

Search for lepton-flavour violating decays of the Higgs boson at the LHC

Dissertation zur Erlangung des Doktorgrades
an der Fakultät für
Mathematik, Informatik und Naturwissenschaften
Fachbereich Physik
der Universität Hamburg

vorgelegt von
Annika Vanhoefer

Hamburg
2017

Gutachter der Dissertation: Prof. Dr. Peter Schleper
Prof. Dr. Elisabetta Gallo

Gutachter der Disputation: Prof. Dr. Peter Schleper
Prof. Dr. Elisabetta Gallo
Prof. Dr. Caren Hagner (Vorsitz)
Prof. Dr. Gudrid Moortgat-Pick
Dr. Alexei Raspereza

Datum der Disputation: 08.05.2017

Abstract

In this thesis, a search for lepton-flavour violating (LFV) decays of the Higgs boson to an electron and a tau lepton, denoted as $e\tau$, is presented. LFV decays of the Higgs boson are forbidden in the Standard Model of particle physics (SM). However, extensions of the SM, such as multi-Higgs doublet models, can introduce LFV Yukawa couplings. The search is performed for tau decays to a muon and corresponding neutrinos, denoted as $e\tau_\mu$. Proton-proton collision data were collected with the CMS detector in 2012 and 2015, while the LHC was operating at a centre-of-mass energy of $\sqrt{s} = 8$ TeV and $\sqrt{s} = 13$ TeV, respectively, accounting for integrated luminosities of 19.7 fb^{-1} and 2.3 fb^{-1} . The dominant SM background contributions for this search are coming from the Drell-Yan process, misidentified leptons, and the top-quark pair-production process. Top-quarks predominantly decay into a b quark and a W boson. Jets originating from b quarks can be identified, with processes including c-quark jets being the main background. The impact of heavy flavour modelling on the misidentification of c-quark jets as b-quark jets is studied. Differences in the charmed-hadron modelling regarding the production fraction, decay, and momentum fractions affect the discrimination. Furthermore, the misidentification probability is found to depend on the type of the charmed-hadron. In the LFV search, a new algorithm is developed to improve the optimisation of the event selection, which determines iteratively the best set of requirements on the discriminating variables. The search is done in three exclusive jet categories. Improvements on the modelling on the misidentified lepton background are studied.

No excess of events is observed in the reconstructed mass distribution of the Higgs boson. An observed (expected) 95% CL upper limit on the branching fraction $B(H \rightarrow e\tau) < 0.74\%$ ($< 0.85^{+0.37}_{-0.25}\%$) can be set using the collision data at $\sqrt{s} = 8$ TeV. Results were combined with the channel for hadronic tau decays and an observed (expected) 95% CL upper limit on the branching fraction $B(H \rightarrow e\tau) < 0.69\%$ ($< 0.75\%$) was obtained, with the $e\tau_\mu$ channel being the most sensitive one. This model independent limit improves upon indirect limits from low energy measurements. An observed (expected) 95% CL upper limit on the branching fraction $B(H \rightarrow e\tau) < 2.01\%$ ($< 1.32^{+0.59}_{-0.39}\%$) can be set using the collision data at $\sqrt{s} = 13$ TeV. This limit is weaker due to smaller luminosity. However, the sensitivity was improved with respect to the 8 TeV analysis, which will be important for the analysis of the 2016 collision data with higher integrated luminosity.

Kurzfassung

In dieser Arbeit wird eine Suche nach Lepton-Flavour verletzenden (LFV) Zerfällen des Higgs-Bosons in ein Elektron und ein τ -Lepton vorgestellt. Solche Zerfälle sind im Standardmodell (SM) der Teilchenphysik nicht erlaubt. Erweiterungen des SM, wie Modelle mit mehreren Higgs-Doublets, sagen LFV Yukawakopplungen vorher. Die Suche wird für τ -Leptonen durchgeführt, welche in ein Myon und entsprechende Neutrinos zerfallen. Verwendet wurden Daten aus Proton-Proton-Kollisionen, welche in den Jahren 2012 und 2015 am LHC mit einer integrierten Luminosität von 19.7 fb^{-1} bzw. 2.3 fb^{-1} und mit einer Schwerpunktsenergie von $\sqrt{s} = 8 \text{ TeV}$ bzw. $\sqrt{s} = 13 \text{ TeV}$ mit dem CMS-Detektor genommen wurden. Der Drell-Yan-Prozess, fehlidentifizierte Leptonen und die Top-Quark-Paarproduktion sind die Hauptuntergründe dieser Suche. Top-Quarks zerfallen fast immer in ein W-Boson und ein b-Quark, dessen Jet identifiziert werden kann, wobei Jets aus c-Quarks den Hauptuntergrund bilden. Der Einfluss der Modellierung dieser Hadronen auf die Fehlidentifizierung als b-Quark-Jets wurde untersucht. Unterschiede im Bezug auf Produktion, Zerfall und Impulsanteil dieser Hadronen haben einen Einfluss auf die Identifizierung. Außerdem hängt die Wahrscheinlichkeit für eine Fehlidentifizierung vom Typ des Hadrons ab. Für die Ereignisauswahl der LFV-Suche wurde ein neuer Optimierungsalgorithmus entwickelt, der iterativ die beste Selektion der Ereignisse anhand der diskriminierenden Variablen ermittelt. Die Suche ist in drei exklusive Jet-Kategorien eingeteilt. Es wurden Möglichkeiten untersucht, die Modellierung des fehlidentifizierten Leptonuntergrunds zu verbessern.

Es wurde kein Überschuss an Ereignissen in der Massenverteilung des Higgs-Bosons beobachtet. Die beobachtete (erwartete) 95% CL Obergrenze auf das Verzweigungsverhältnis beträgt $B(H \rightarrow e\tau) < 0.74\%$ ($< 0.85^{+0.37}_{-0.25}\%$) für die 8 TeV-Daten. Für deren Kombination mit dem hadronischen τ -Lepton-Zerfall Kanal beträgt sie $B(H \rightarrow e\tau) < 0.69\%$ ($< 0.75\%$), wobei der $e\tau_\mu$ -Kanal am sensitivsten ist. Diese modellunabhängigen Ausschlussgrenzen sind besser als indirekte Ausschlussgrenzen von Niederenergiemessungen. Die Ausschlussgrenze der 13 TeV-Daten beträgt $B(H \rightarrow e\tau) < 2.01\%$ ($< 1.32^{+0.59}_{-0.39}\%$), aufgrund der geringeren Luminosität. Jedoch wurde die Sensitivität verbessert, was wichtig für die Analyse der Daten ist, welche 2016 genommen wurden und eine höhere integrierte Luminosität besitzen.

Contents

1. Introduction	1
2. The Standard Model	5
2.1. Phenomenological overview	5
2.1.1. Quantum electrodynamics	6
2.1.2. Electroweak interaction	7
2.1.3. Quantum chromodynamics	8
2.1.4. The Higgs boson	10
2.2. Overview of theoretical building blocks	10
2.3. The Standard Model	14
2.4. Electroweak symmetry breaking and the Higgs boson	16
2.5. Main shortcomings of the Standard Model	18
3. Lepton-flavour violating decays of the Higgs boson	19
3.1. Introduction of lepton-flavour violating decays	20
3.2. Constraints from low-energy measurements	25
3.3. Lepton-flavour violating decays of the Higgs boson at the LHC	31
4. Experimental Setup	37
4.1. The Large Hadron Collider	37
4.2. The Compact Muon Solenoid experiment	43
4.2.1. The coordinate system of CMS	44
4.2.2. Kinematic quantities	44
4.2.3. Detector requirements and layout	45
4.2.4. Magnet	48
4.2.5. Inner Tracking system	48
4.2.6. Electromagnetic Calorimeter	49
4.2.7. Hadron Calorimeter	50
4.2.8. Muon System	52
4.2.9. Trigger and data acquisition system	53
4.3. Luminosity measurement	55
5. Event simulation	59
5.1. Short distance simulation	61
5.2. Hadronisation models	64
5.3. Soft hadron-hadron physics modelling	67
5.4. Parameters and tuning	67
5.5. Monte Carlo generators	68
5.6. Detector simulation	74

6. Event reconstruction	77
6.1. Track and primary-vertex reconstruction	79
6.1.1. Hit reconstruction	79
6.1.2. Track reconstruction	80
6.1.3. Special track reconstruction	82
6.1.4. Primary-vertex reconstruction	83
6.1.5. Beam spot reconstruction	84
6.2. Electron reconstruction	85
6.2.1. Electron reconstruction	85
6.2.2. Electron charge estimation	90
6.2.3. Electron momentum estimation	90
6.2.4. Electron identification	92
6.2.5. Electron isolation	93
6.3. Muon reconstruction and identification	95
6.3.1. Muon reconstruction	95
6.3.2. Muon momentum assignment	96
6.3.3. Muon sources (in simulation)	97
6.3.4. Identification	98
6.3.5. Muon isolation	100
6.4. Lepton efficiency estimation using a tag-and-probe technique	101
6.4.1. Tag-and-probe technique	101
6.4.2. Electron efficiencies	102
6.4.3. Muon efficiencies	102
6.5. Particle-flow event reconstruction	104
6.5.1. Fundamental elements	104
6.5.2. Link algorithm	105
6.5.3. Particle reconstruction and identification	106
6.5.4. Cluster calibration	108
6.6. Jet reconstruction	109
6.6.1. Jet reconstruction	109
6.6.2. The anti- k_t algorithm	110
6.6.3. Jet energy scale and resolution	112
6.7. Missing transverse energy	117
6.7.1. Reconstruction of missing transverse energy	117
6.7.2. Corrections for reconstructed missing transverse energy	118
6.7.3. Scale and resolution of missing transverse energy	120
6.8. Identification of b-quark jets	121
6.8.1. Reconstructed objects used in b-quark jet identification	122
6.8.2. The Combined Secondary Vertex (CSV) algorithm	124
6.8.3. Efficiency of the Combined Secondary Vertex (CSV) algorithm	125
6.8.4. Identification of b-quark jets in LHC Run II	130

7. Impact of heavy flavour modelling on the misidentification of c-quark jets	135
7.1. Simulation, selection and categorisation of jets	136
7.2. Differences in heavy flavour modelling	139
7.2.1. Charmed hadron production fractions	139
7.2.2. Charmed hadron decays	140
7.2.3. Charmed hadron momentum fraction	142
7.2.4. Misidentification probability of c-quark jets	144
7.2.5. Conclusions	145
7.3. Reweighting HERWIG to PYTHIA values	146
7.3.1. Reweighting one quantity	146
7.3.2. Combined reweighting	150
7.4. Reweighting to the measured values	152
7.4.1. Charmed hadron production fractions	152
7.4.2. Charmed hadron decays	155
7.4.3. Charmed hadron momentum fraction	160
7.4.4. Remaining differences between HERWIG and PYTHIA	166
7.4.5. Effect on the misidentification probabilities of c-quark jets	169
7.5. Systematic uncertainties on misidentification probability of c-quark jets . .	171
7.6. Conclusions	174
8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$	177
8.1. Signal process	178
8.2. Mass reconstruction	181
8.3. Background processes	184
8.4. Preselection of events	187
8.4.1. Recorded data sample	187
8.4.2. Physics objects	188
8.4.3. Corrections	192
8.4.4. Preselection	198
8.5. Background modelling	198
8.5.1. $Z \rightarrow \tau\tau$ background	199
8.5.2. Misidentified lepton background	200
8.6. Background normalisation	213
8.6.1. Drell-Yan: $Z \rightarrow ll$	213
8.6.2. Top background	215
8.6.3. Diboson background	216
8.6.4. W+jets	222
8.7. Profile likelihood method	223
8.8. Systematic uncertainties	223
8.8.1. Normalisation uncertainties	223
8.8.2. Normalisation uncertainties of the background processes	224
8.8.3. Normalisation uncertainties of the Higgs boson processes	225
8.8.4. Shape uncertainties of the collinear mass	226
8.9. Discriminating variables	227
8.9.1. Collinear mass distributions after preselection	227
8.9.2. Discriminating variables used in the $\sqrt{s} = 8$ TeV analysis	229
8.9.3. New discriminating variables	235

8.10. Search for lepton-flavour violating Higgs-boson decays at $\sqrt{s} = 8$ TeV . . .	237
8.11. Optimisation of the event selection	246
8.11.1. Background modelling	246
8.11.2. Optimisation method	247
8.11.3. Optimisation for maximising the significance	250
8.11.4. Optimisation for minimising the expected 95% CL upper limit . . .	252
8.11.5. Optimisation for expected limit using new discriminating variables .	254
8.12. Expected 95% CL upper limits on the branching fraction $B(H \rightarrow e\tau_\mu)$. . .	255
8.13. Event selection	264
8.14. Results	267
9. Summary and discussion of results	277
9.1. Comparison to other constraints	278
9.2. Possible improvements	282
Appendices	283
A. Mass resolution	284
A.1. Derivation of the collinear mass M_{col} reconstruction for $H \rightarrow e\tau_\mu$	284
B. Event simulation	289
C. Misidentified lepton background	292
D. Control regions	299
E. Search for LFV decays of the Higgs boson at $\sqrt{s} = 8$ TeV	304
F. Discriminating variables of the 8 TeV analysis at $\sqrt{s} = 13$ TeV	306
G. Expression for the 95% CL upper limit on the branching fraction	312
H. Results of the search at $\sqrt{s} = 13$ TeV	314
I. New luminosity estimate for LHC Run II	319
I.1. Impact on the search for LFV decays of the Higgs boson	320
I.2. Results	326
J. Search for LFV decays of the Higgs boson with the ATLAS experiment	328
List of Figures	330
List of Tables	336
Bibliography	340
List of Publications	354

1. Introduction

Particle physics explores the elementary particles and fundamental forces between them in order to understand the nature of the universe. Knowledge about the universe is gained by experimental tests of explanations and predictions. At present, the Standard Model (SM) of particle physics is consistent with all known particles and all known interactions, but gravity. Forces are mediated by the exchange of particles, referred to as *gauge bosons*. The Higgs mechanism explains how by introduction of a new scalar field the gauge bosons acquire mass. At the same time, this requires a new particle, the *Higgs boson*. Fermions acquire mass through interacting with the Higgs field via *Yukawa interactions*. In the SM, an eigenstate basis can be chosen so that both, the mass matrix and the Yukawa interaction matrix, are diagonalised simultaneously. Thus, there are no decays of the Higgs boson in the SM that violate lepton-flavour conservation.

The SM has been successfully tested over a wide range of energies. Despite its success, the SM is considered as an effective theory as it fails to include gravity. Furthermore, it cannot describe observations such as *Dark Matter*, which points to the existence of new particles. Another observation which cannot be explained by the SM is the accelerated extension of the universe which could be explained by the introduction of *Dark Energy*.

The understanding of the flavour structure of fermions is still ongoing as the SM cannot explain why there are six quarks and six leptons which can be arranged into three generations. There are three replications for each of the four fermion types which only differ by their mass. It is not known why there are three generations. The mass values of the SM particles, including the mass of the Higgs boson, cannot be derived within the SM.

CP-violating effects of the SM cannot explain the large difference between matter and anti-matter in the universe. The only observed source of CP-violation is the Cabibbo-Kobayashi-Maskawa (CKM) matrix, a 3×3 unitary matrix which can be parameterised by three mixing angles and one complex CP-violating phase [1]. The mass eigenstates are not equal to the weak eigenstates and the CKM matrix describes the mixing of the quarks.

Neutrinos are strictly massless within the original formulation of the SM because there is no Dirac mass term due to the absence of right-handed neutrinos and there is no Majorana mass term due to exact conservation of baryon and lepton number. However, neutrino experiments have established oscillations of neutrinos caused by nonzero neutrino masses. Neutrino mixing is described by the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix [1].

1. Introduction

The PMNS matrix is determined by three mixing angles and a phase which could lead to additional CP-violation. Furthermore, there is a remarkable difference between the magnitudes of the very small neutrino masses ($m < 2$ eV [1]) and the masses of the charged leptons and quarks (~ 173 GeV for the top quark [1]) which cannot be explained by the SM. An open question is if there is any relation between the quark mixing and the neutrino mixing. Gauge invariance cannot constrain the flavour structure of the Yukawa interactions within the SM. New physics could explain some of the symmetries and open questions. Grand Unification theories (GUTs) try to unify all known interactions but gravity in a single gauge group, while extra dimensions could link flavour to the geometry of this extra dimensions [2]. A simple extension of the Higgs sector with at least two Higgs doublets could add additional CP-violation. Multi-Higgs doublet models can have tree level Higgs-mediated flavour changing neutral currents leading to lepton-flavour violating (LFV) Yukawa couplings [2].

To extend the energy range covered by experiments, the Large Hadron Collider (LHC) was built. It was designed to run at centre-of-mass energy of $\sqrt{s} = 14$ TeV. It is the most powerful particle collider to date and was built to discover the Higgs boson, to measure its properties, to search for new physics and to test predictions of theories for extensions of the SM. Two general-purpose detectors were built by the ATLAS and CMS collaborations. In 2012, a Higgs boson with a mass of 125 GeV was discovered by the ATLAS and CMS collaborations [3, 4]. Measurements of its properties and precision measurements have to decide if it is the SM Higgs boson or if it is a Higgs boson of a possible extension of the SM. Thus, properties such as production and the decay of the Higgs boson and its spin are being measured. No significant deviations within the uncertainties from the SM Higgs boson have been observed so far.

In this thesis, a search for lepton-flavour violating (LFV) decays of the Higgs boson to an electron and a tau which decays further to a muon and neutrinos, referred to as $e\tau_\mu$, is presented. This search is performed with proton-proton collision data collected at a centre-of-mass energy of $\sqrt{s} = 8$ TeV during 2012 and at $\sqrt{s} = 13$ TeV during 2015 corresponding to an integrated luminosity of 19.7 fb^{-1} and 2.3 fb^{-1} , respectively. Results of the search done at $\sqrt{s} = 8$ TeV are published in [5]. Similar methods for the background processes and similar discriminating variables were used for the 8 TeV and the 13 TeV analysis, but some improvements have been made. New algorithms were developed in order to optimise the requirements on discriminating variables for defining an event selection, optimising for significance or for expected limit for the 8 TeV analysis and the 13 TeV analysis, respectively. The search is done in three exclusive jet categories: 0-jet category, 1-jet category, and the 2-jet category.

This thesis is structured as follows. An overview of the Standard Model of particle physics is given in **Chapter 2**. A short review on lepton-flavour violating decays of the Higgs bosons is given in **Chapter 3**. It is described how lepton-flavour violating Yukawa-couplings can be introduced by new physics and how they are constrained by low-energy measurements [6]. Then, the experimental setup is discussed in **Chapter 4**. The Large Hadron Collider (LHC) [7] and its experiments are introduced with the focus on the Compact Muon Solenoid (CMS) experiment. After it is explained how events are recorded by the CMS detector, trigger and data acquisition system, an introduction in the event simulation is given in **Chapter 5**. An overview of the event reconstruction is given in **Chapter 6**. The study on the impact of heavy flavour modelling on the misidentification of c-quark jets as b-quark jets is presented in **Chapter 7**. It was done as part of technical work for the CMS collaboration. In **Chapter 8**, the search for lepton-flavour violating (LFV) decays of the Higgs boson to an electron and a tau which further decays to a muon and neutrinos is presented. A summary of this thesis is given in **Chapter 9**. Results of the ATLAS and CMS collaboration on LFV decays of the Higgs boson are discussed and compared to the results presented in this thesis including the search done at $\sqrt{s} = 8$ TeV.

2. The Standard Model

The Standard Model (SM) of particle physics describes the constituents of matter and their interactions at the most fundamental level. It is a renormalisable quantum field theory with a $SU(3) \times SU(2) \times U(1)$ symmetry structure. All known interactions, but gravity, are included in the SM.

This chapter follows partly the corresponding chapter in my master thesis [8], which is based on [9]. However, a stronger focus is set on the mathematical basis of the SM in this thesis. A complete introduction into the mathematical basis of the SM can be found in [10]. First, a phenomenological and a theoretical overview of the elementary particles, the fundamental interactions and the Higgs boson are given [1, 8–10], before the Standard Model is introduced [8–10]. Then, it is discussed how weak gauge bosons and fermions acquire mass via the Higgs mechanism which breaks the electroweak symmetry [1, 10]. Main shortcomings of the SM are listed at the end of the chapter.

2.1. Phenomenological overview

Elementary particles can be classified according to their intrinsic angular momentum, their *spin*. There are two different types of particles, *fermions* which are particles with half-integer spin ($\frac{1}{2}, \frac{3}{2}, \frac{5}{2}, \dots$), and *bosons* which are particles with integer spin ($0, 1, 2, \dots$).

All matter consist of *quarks* and *leptons*, which are elementary fermions with half-integer spin. There are six quark flavours: up (u), down (d), strange (s), charm (c), bottom (b), and top (t). There are three different types of charged leptons: the electron (e), the muon (μ), and the tau τ . Each lepton has its corresponding neutral partner: the electron neutrino (ν_e), the muon neutrino (ν_μ), and the tau neutrino (ν_τ). The six flavours of leptons and quarks can be arranged into three *generations*, respectively. Each fermion has an anti-particle with the same properties, but opposite charges. Figure 2.1 shows gives a summary of the elementary particles of the SM.

The fundamental interactions are: gravitation, electromagnetism, weak interaction, and strong interaction. Gravitation is the weakest force and it is described by the general theory of relativity. It has an infinite range and it is important for macroscopic objects and the large-scale structure of the universe. Among the other forces, only electromagnetism

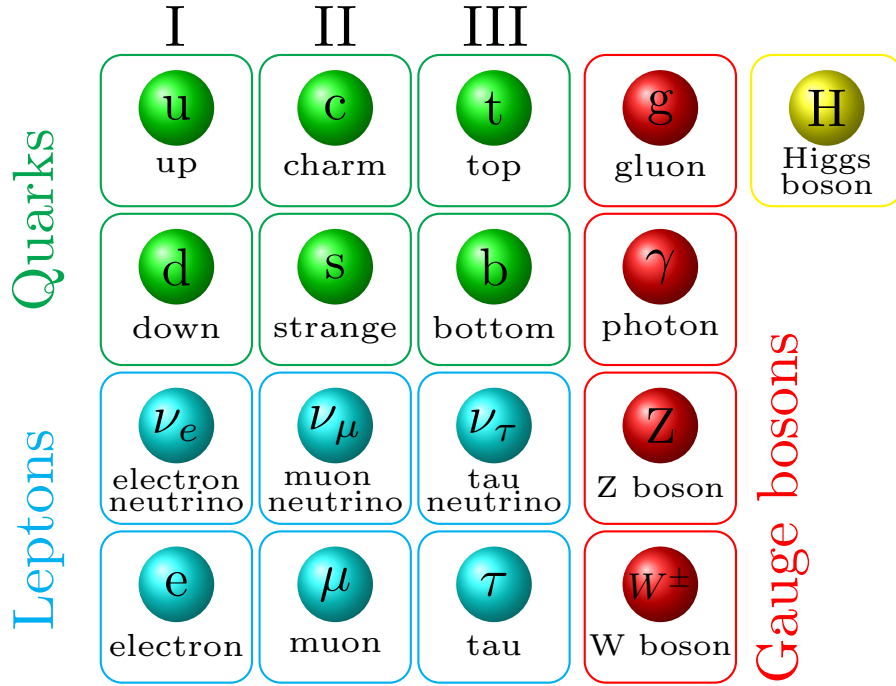


Figure 2.1.: Elementary particles of the Standard Model

has an infinite range, though positive and negative charges cancel each other on average at larger scales. All interactions, but gravitation, are part of the Standard Model and can be described as quantum fields and their interactions are mediated by field quanta, the *gauge bosons*, which have spin 1. Quantum electrodynamics (QED) is the relativistic quantum field theory of electrodynamics. The weak interaction and electromagnetism can be unified to one theory, the electroweak theory. Quantum chromodynamics (QCD) is the theory of the strong interaction. In the following, a short phenomenological overview of these interactions is given.

2.1.1. Quantum electrodynamics

Quantum electrodynamics (QED) is the quantum field theory of electromagnetism. Its mediator is the photon (γ), which acts between electrically charged particles. Charged leptons have a charge of ± 1 , with particles having a negative charge. Neutrinos are electrically neutral, thus they do not interact via the electromagnetic interaction. Up-type quarks have a charge of $+2/3$, while down-type quarks have a charge of $-1/3$. The photon itself carries no electrical charge. It is massless, therefore the electromagnetic interaction has an infinite range. In QED, the elementary process is the emission and absorption of a photon: $e^- \rightarrow e^- + \gamma$, which is illustrated in Figure 2.2.

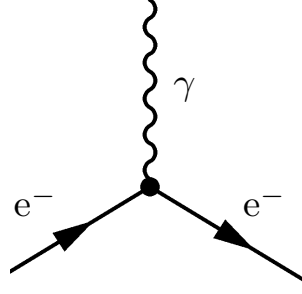


Figure 2.2.: Fundamental QED vertex – Emission or absorption of a photon by electrons (electrically charged particles).

All processes of the electromagnetic interaction can be described by combining two or more of the fundamental vertices to *Feynman diagrams*, which represent mathematical expressions for calculating the probability amplitudes for a given process. Antiparticles are indicated as arrows going backwards in time. Initial and final-state particles define the physical process. Internal lines represent particles, which cannot be observed. The total probability amplitude of a process is proportional to the squared sum of all Feynman diagrams which represent the same process, i.e. which have the same initial and final state. These amplitudes can have infinite contributions. Finding the infinite parts and separating them from the finite parts of the amplitude is the aim of *renormalisation*. In the regularisation step, the infinite observable is modified by adding an additional parameter in order to make it finite. The resulting observable depends on the additional parameter, but it can be computed without divergences. Then, the end result is obtained by taking the parameter to its physical limit. As real quantum systems can get very complex and as it is very difficult to find exact solutions, perturbation theory plays an important part in the calculations of quantum field theory, described in Section 2.2.

The strength of the electromagnetic interaction between electrically charged particles is characterised by the fine-structure constant α . Each vertex introduces a factor proportional to $\sqrt{\alpha}$, which depends on the energy scale. Virtual e^+e^- pairs can spontaneously be produced around a "bare" electron reducing the observed charge at larger distances. This is referred to as *charge screening*. The value of α decreases with energy to a constant value of $\approx 1/137$ for long-distance interactions. Very precise QED predictions can be obtained using only lower order diagrams, as the higher-order diagrams can be neglected due to the small value of α .

2.1.2. Electroweak interaction

The weak interaction was first observed in certain types of radioactive decays. It is the only interaction that can change the flavour of particles. Leptons and quarks can interact weakly. Neutrinos can only interact via the weak interaction. The interaction is mediated

2. The Standard Model

by three gauge bosons: the W ($W^\pm$) and Z bosons (Z^0). Figure 2.3 shows the fundamental vertices of the weak interaction. The Z boson can couple to two fermions of same flavour, while the W boson couples to fermions of different flavour. Furthermore, both bosons can interact with each other. As the W bosons are electrically charged, they can additionally couple to photons. The weak interaction only acts over very short distances in the order of 10^{-18} m, because the Z and W bosons have large masses, about 91 GeV and 80 GeV [1], respectively.

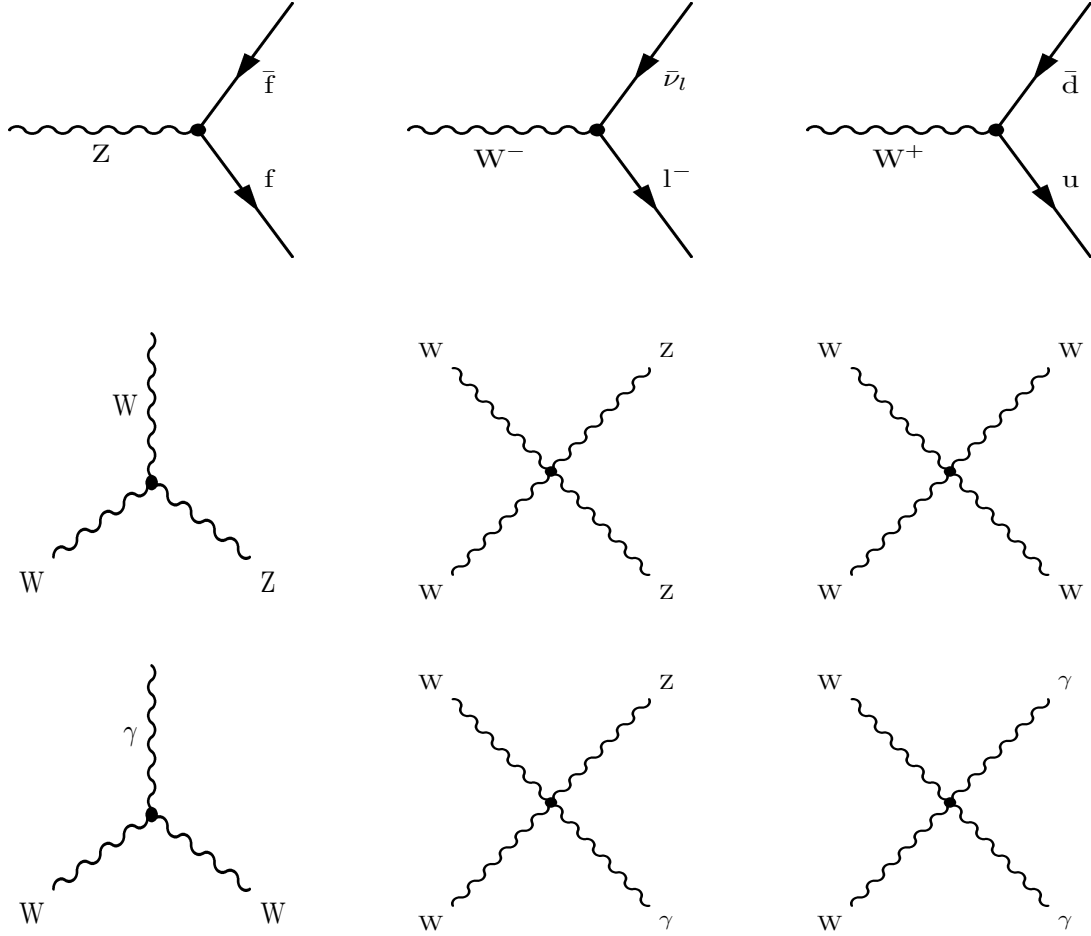


Figure 2.3.: Fundamental vertices of the weak interaction. Fermions are denoted f and their anti-particles are denoted $\bar{f}$.

2.1.3. Quantum chromodynamics

Quantum chromodynamics (QCD) is the quantum field theory of the strong interaction, which binds neutrons and protons into an atomic nuclei. The strong interaction has a short range in the order of 10^{-15} m, the size of the proton. It is mediated by a massless gauge boson called *gluon* (g) which acts on particles with a colour charge. There are three of such charges which are denoted red, green, and blue (r, g, b). The corresponding

anti-colours are denoted $\bar{r}$, $\bar{g}$, and $\bar{b}$. Quarks are the only fermions which carry a colour charge. A fundamental process of the strong interaction is the process, where a quark emits or absorbs a gluon: $q \rightarrow q + g$. The gluons carry colour themselves, a colour and an anti-colour, therefore they interact with each other. There are eight types of gluons. Figure 2.4 shows the fundamental vertices of the strong interaction.

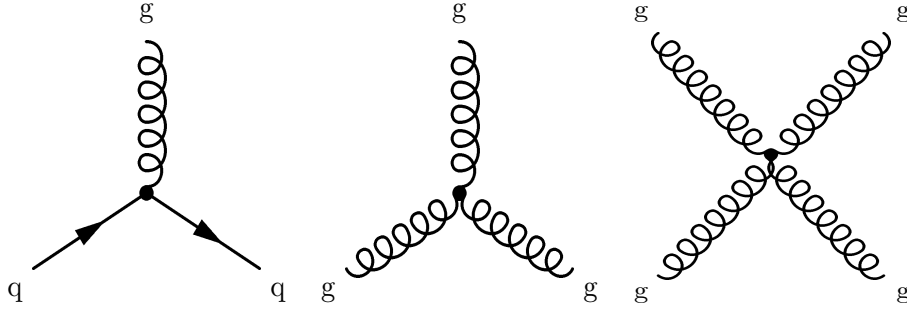


Figure 2.4.: Fundamental vertices of the strong interaction.

A bare colour charge is surrounded by a "sea" of virtual quarks and gluons. At shorter distances, corresponding to higher energies, the observed charge lessens until only the bare charge is seen. The phenomenon that the interaction strength gets weaker with increasing energy and decreasing distance is referred to as *asymptotic freedom*. The observed charge increases with the distance, resulting in a strong attractive force between two distant charges. At large distance, the potential energy between two quarks is large enough to create a real quark-antiquark pair from the vacuum. This process is known as *fragmentation* or *hadronisation*. Two separating quarks always hadronise to colourless particles. The observation that only colourless bound states have been observed in nature is referred to as *colour confinement*. The strong interaction binds quarks into composite states called *hadrons*. *Baryons* are composed of three quarks with spin $1/2$ or $3/2$. Their counterparts are made from anti-quarks. *Mesons* are composed of a quark and an anti-quark. They are spin-0 or spin-1 particles. Thus, mesons are bosons, while baryons are fermions.

The coupling strength of the strong interaction α_S describes the dependence of the effective charge on the distance between them. In the lowest order, $\alpha_S(Q)$ is given by:

$$\alpha_S = \frac{6\pi}{(33 - 2n_f) \ln(Q/\Lambda_{\text{QCD}})}, \quad (2.1)$$

where Q is the momentum transfer in a given process, n_f the number of flavours which can participate in the process, and Λ_{QCD} corresponds to the energy boundary of hadronisation [10]. For $Q \approx \Lambda_{\text{QCD}}$, quarks and gluons interact strongly and form hadrons, while $\alpha_S(Q)$ becomes smaller for $Q \gg \Lambda_{\text{QCD}}$ and quarks and gluons interact with each other only weakly. At low energies, perturbative methods cannot be used to investigate the theory due to the large coupling strength.

2.1.4. The Higgs boson

In the SM, the gauge bosons are required to have no mass, but non-zero masses were measured for the W and Z bosons. They receive their mass by interacting with the Higgs field and via the *Higgs mechanism*. Also the fermions receive mass via this mechanism. The Higgs boson (H) is the field quanta of the Higgs field and has spin-0. Figure 2.5 shows the fundamental vertices of the Higgs boson with the fermions, the gauge bosons (V) of the weak interaction and its self-interactions. It couples to two fermions of the same type, two W or Z bosons, and with further Higgs bosons.

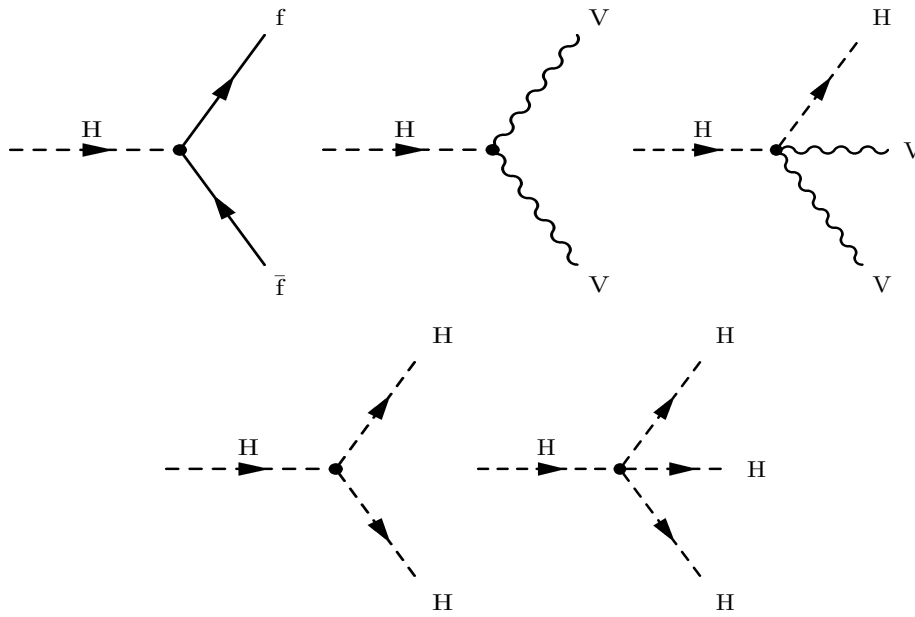


Figure 2.5.: Fundamental vertices of the Higgs boson. Fermions are denoted f and anti-fermions are denoted $\bar{f}$. Gauge bosons of the weak interaction are denoted V .

2.2. Overview of theoretical building blocks

The Standard Model is a renormalisable quantum field theory. It has a $SU(3) \times SU(2) \times U(1)$ symmetry structure. Understanding the general, abstract details of structures is the aim of group theory. The dynamics of a system is described by the corresponding Lagrangian of the theory. A type of quantum field theories are gauge theories, where gauge fields are introduced by requiring symmetry under local transformations. Finally, the idea of spontaneous symmetry breaking is discussed. A complete introduction into mathematical basis of the Standard Model, including group theory, can be found in [10].

Group theory: A *group* $(G, \star)$ is a set of objects G and some operations $\star$ on those objects. A set of elements g_i is a group under some operation if the group axioms are satisfied: closure, associativity, existence of an identity element, and an existence of an inverse element. A group is *Abelian* if the elements of the group commute under the group operation ($g_i \star g_j = g_j \star g_i$). Groups which are not Abelian are called *non-Abelian*. The *order* n of the group is given by the number of elements in the group. A group can be represented by $n \times n$ matrices.

A special type of groups are *Lie groups*, which are parameterised by one or more continuous variables. Transformations of elements $\in \mathbb{C}$ can be described as matrices R . These matrices are unitary if the Hermitian conjugate is also the inverse of the transformation, $R^\dagger = (R^*)^T = R^{-1}$. A set of such $n \times n$ matrices form the *unitary group* $U(n)$, which preserves the norm under transformations, and thus the probability amplitude. Some of these unitary matrices R have a determinant $\det |R| = 1$. The volume is preserved under such transformations. A set of such matrices form the *special unitary group* $SU(n)$.

Lagrangian formalism The information about the dynamics of a system is given by its Lagrangian as function of generalised coordinates q and their time derivatives $\dot{q}$. In classical physics, the Lagrangian is defined as the kinetic energy T minus the potential energy V :

$$L = T(q, \dot{q}) - V(q). \quad (2.2)$$

Fields are functions of spacetime. In field theory, the Lagrangian is replaced by the Lagrangian density $\mathcal{L}$, which is a function of the fields $\phi(x^\mu)$ in the system and their derivatives $\partial_\mu \phi$. The *action* S of the system is defined as:

$$S = \int d^4x \mathcal{L}. \quad (2.3)$$

Quantum field theory Quantum field theory (QFT) is a relativistic quantum theory of interactions. Particles are interpreted as fields and forces between particles are mediated by other particles, the gauge bosons. In quantum mechanics, space is treated as an operator, while time is treated as in classical physics, as a parameter. However, space and time have to be treated in the same way for a relativistic theory. This is solved in QFT by quantising in a new way. The fields are reinterpreted no longer as states, but as operators, which are parameterised by the spacetime coordinates. Commutation relations are imposed on the field operators. It is not possible to measure both the field and its time rate of change simultaneously to infinite precision at the same spacetime point. The state with lowest possible energy level is called *vacuum*. Acting on the vacuum with an field operator produces a state with some energy. The Lorentz transformations known from special relativity are elements of the Lorentz group. This group contains two

2. The Standard Model

copies of the SU(2) group, where the SU(2) group represents quantum mechanical spin. There are different representations, which are characterised by the value j , the spin of the particle. Consequently, a representation (j, j') of the Lorentz group is characterised by two values j and j' for each SU(2) subgroup. There are three physically acceptable irreducible representations of the Lorentz group:

- the $(0, 0)$ scalar representation acting on *scalars*
- the $(\frac{1}{2}, 0) \oplus (0, \frac{1}{2})$ left-handed/right-handed spinor representations acting on *spinors*
- the $(\frac{1}{2}, \frac{1}{2})$ vector representation acting on *vectors*.

There are two types of spinors, left-handed Ψ_L and right-handed Ψ_R , which transform the same under rotations, but they transform differently under boosts. Some fundamental properties of these objects are outlined in the following.

Scalar fields ϕ are spin-0 fields. Previously, spin was defined as intrinsic angular momentum. *Spin* is not a rotation in spacetime, but in the spinor space of SU(2). Complex scalar fields have the form $\phi = \phi_1 + i\phi_2$, with two real degrees of freedom. The fields ϕ and $\phi^\dagger$ are treated as independent fields and the Lagrangian density is given as function of both fields and their derivatives.

Spinor fields Ψ are spin-1/2 fields and consist of four complex fields Ψ_i :

$$\Psi = \begin{pmatrix} \Psi_1 \\ \Psi_2 \\ \Psi_3 \\ \Psi_4 \end{pmatrix}, \quad \Psi^\dagger = (\Psi_1^*, \Psi_2^*, \Psi_3^*, \Psi_4^*). \quad (2.4)$$

To have Lorentz invariant terms in the Lagrangian density, the adjoint spinor $\bar{\Psi} = \Psi^\dagger \gamma^0$ has to be defined. The γ^μ matrices satisfy the anti-commutation relation (Clifford algebra):

$$\{\gamma^\mu, \gamma^\nu\} = \gamma^\mu \gamma^\nu + \gamma^\nu \gamma^\mu = -2\eta^{\mu\nu} \mathbb{1}, \quad (2.5)$$

with Minkowski metric $\eta^{\mu\nu}$. The Lagrangian density is given as function of the fields Ψ , $\bar{\Psi}$, and their derivatives. Parity transformations Λ_P are one type of Lorentz transformations, which switch the handedness of a coordinate frame. Left-handed spinors are transformed into right-handed spinors under parity transformations. To be Lorentz invariant, spinor fields Ψ should have a left-handed and a right-handed component, with the adjoint field of a left-handed spinor field $\bar{\Psi}_L$ being the same field as the right-handed spinor field Ψ_R :

$$\bar{\Psi}_L = \Psi_R, \quad \bar{\Psi}_R = \Psi_L. \quad (2.6)$$

The projection operators $P_{\pm}$ are the projection operators onto the left- and right-handed parts of the field Ψ :

$$P_+ \Psi = \Psi_R, \quad P_- \Psi = \Psi_L. \quad (2.7)$$

Another important transformation is the *charge conjugation*: $\Psi \rightarrow C\bar{\Psi}^T = -i\gamma^2\Psi^*$, which swaps the charge of the field.

Vector fields A_{μ} have spin 1. They have four components that transform as a spacetime vector under Lorentz transformations. Lorentz invariant terms are:

$$A_{\mu}A^{\mu}, \quad (\partial_{\mu}A_{\nu})(\partial^{\mu}A^{\nu}), \quad (\partial_{\mu}A^{\mu})(\partial_{\nu}A^{\nu}). \quad (2.8)$$

In summary, the fundamental fermions are represented as spinor fields, while gauge bosons are represented as vector fields in QFT. Combining the Lagrangian densities of both types of fields leads to two fields, which do not interact with each other. An *interaction term* has to be added which mixes both fields. Such a term can be introduced using *gauge theory*.

Gauge theory A type of field theories are *gauge theories*, which are theories in which the Lagrangian is invariant under a continuous group of local transformations. For example, a spinor field can be transformed with the *global* $U(1)$ transformation $\Psi \rightarrow e^{i\alpha}\Psi$, with α being a real parameter. The transformation is just a phase and acts on the field in exactly the same way at every point in spacetime. A *local* transformation has a spacetime dependence. In this example of the $U(1)$ transformation, α would depend on spacetime: $\alpha = \alpha(x^{\mu})$. The possible transformations of the same type are associated with a Lie group, the *symmetry group* of the theory. All elements of a group can be expressed in terms of the *group generators*. First, a Lagrangian density is defined for a fermion field with an equation of motion that satisfies the Dirac equation. This Lagrangian density is not invariant under a local gauge transformation. A new vector field A_{μ} has to be introduced, called *gauge field*, for each group generator T_a to ensure invariance under the local group transformations, called *gauge invariance*. The vector field can be expressed in terms of the group generators, $A_{\mu} = A_{\mu}^a T_a$, and the *gauge bosons* are the quanta of this fields. The derivative operator ∂_{μ} in the Lagrangian density is replaced by the *covariant derivative* D_{μ} , which is ∂_{μ} with a term proportional to a gauge field. Local transformations can be written in form of $U = e^{ig\chi_a(x^{\mu})T_a}$, where g is a real parameter, the coupling constant, and $\chi_a(x^{\mu})$ is a scalar with a spacetime dependence. The covariant derivative is then given by $D_{\mu} = \mathbb{1}^{n \times n} \partial_{\mu} + igA_{\mu}^a T_a$. Interaction terms are introduced by the second term of the covariant derivative. A kinetic term, proportional to the derivatives of the fields, can be introduced for including dynamics of the vector field. Vector fields with mass break the local gauge symmetry, therefore the vector fields are massless, i.e. have no terms

2. The Standard Model

proportional to $A_\mu A^\mu$. All fundamental interactions of the Standard Model can be derived from this general principle of local gauge invariance. As the SM W and Z bosons are measured to be massive, a mechanism is needed in order that they acquire mass. This mechanism is the *Higgs mechanism* and combines the principles of gauge theory and *spontaneous symmetry breaking*.

Spontaneous symmetry breaking Gauge fields can acquire mass in case of a broken local symmetry. Breaking a global symmetry always results in massless bosons, called *Goldstone bosons* [11,12]. Consider a complex scalar field with a Lagrangian which has a local U(1) symmetry and a potential $V(\phi^\dagger, \phi)$ with the vacuum V_{minimum} given by the circle $|\phi| = v$. Choosing a single ground state breaks the symmetry of the system and it is considered "spontaneous" because there is no external means by which this occurs. The theory can be rewritten in terms of the new vacuum, by a change of basis. The field is expanded around the constant vacuum value, writing the fields in terms of fluctuations around the chosen vacuum.

The *Higgs mechanism* combines the principles of gauge invariance with spontaneous symmetry breaking. For a local U(1) symmetry, the vacuum can be chosen in a way that the vacuum is real and that ϕ is always real, therefore ϕ can be expanded as $\phi = v + h$, with h being a real scalar field. Breaking the local U(1) symmetry results in a real scalar h with a mass and the field A_μ , originating from the local gauge symmetry, acquires mass. The *Higgs mechanism* introduces mass into a theory by breaking the local symmetry. A scalar *Higgs field* h is introduced and the gauge field of the broken symmetry acquires mass. In general, gauge fields corresponding to group generators of broken symmetries acquire mass, while those corresponding to generators of unbroken symmetries do not acquire a mass.

2.3. The Standard Model

The Standard Model is a non-Abelian gauge theory based on the group $SU(3)_C \times SU(2)_L \times U(1)_Y$, where C stands for colour, Y stands for the hypercharge to distinguish $U(1)_Y$ from the $U(1)$ of QED, and L indicates that the interaction only involves left-handed states.

QCD is based on the group $SU(3)_C$, which corresponds to rotations in colour space. The strong interaction is mediated by eight gauge bosons, called gluons. Fermions are described as a triplet of spinors which differ in colour

$$\psi = \begin{pmatrix} \psi_r \\ \psi_b \\ \psi_g \end{pmatrix}. \quad (2.9)$$

Table 2.1.: Overview of the symmetry groups of the Standard Model

Group	Operation	Coupling	Vector field
$U(1)_Y$	phase rotation	g'	B_μ
$SU(2)_L$	rotation in the weak isospin	g	W_μ^a
$SU(3)_C$	rotation in the colour	g_s	G_μ^a

There are eight gauge fields G_μ^a with the gluons being the field quanta, which mediate the strong interaction. Leptons do not interact under the $SU(3)$ colour group, but quarks do.

The electroweak theory is based on the group $SU(2)_L \times U(1)_Y$. Only left-handed fields interact under $SU(2)_L$, therefore only a left-handed neutrino is needed. Left-handed and right-handed fields can interact under the $U(1)_Y$ group, therefore charged leptons have to exist in a left-handed and a right-handed state. Left-handed $SU(2)_L$ doublets

$$\psi_L = \begin{pmatrix} \nu_e \\ e_L \end{pmatrix} \quad (2.10)$$

and right-handed $SU(2)_L$ singlets e_R are introduced. There are 3 gauge fields W_μ^a for $SU(2)_L$, and one gauge field B_μ for $U(1)_Y$. Table 2.1 gives an overview of the symmetry groups the SM is based on. The gauge fields are all massless, but the gauge bosons of the weak theory have mass. Masses can be generated for the W and Z bosons via the Higgs mechanism by breaking the electroweak symmetry. Also quarks and leptons acquire mass via the Higgs mechanism by applying the mechanism to the corresponding fields.

A Higgs field is introduced, which is a complex scalar $SU(2)_L \times U(1)_Y$ doublet field. Breaking the symmetry of the vacuum via the Higgs mechanism results in a broken theory with two separate forces: the broken weak theory and the unbroken theory of electromagnetism associated with the $U(1)_{\text{em}}$ symmetry of QED. The new fields A_μ , $W_\mu^\pm$, and Z_μ^0 are linear combinations of the vector fields of $SU(2)_L \times U(1)_Y$:

$$A_\mu = \sin \theta_W W_\mu^3 + \cos \theta_W B_\mu, \quad (2.11)$$

$$W_\mu^\pm = (W_\mu^1 \mp iW_\mu^2)/\sqrt{2}, \quad (2.12)$$

$$Z_\mu^0 = -\cos \theta_W W_\mu^3 + \sin \theta_W B_\mu, \quad (2.13)$$

with the *weak mixing angle* or Weinberg angle $\theta_W = \tan^{-1} \left(\frac{g'}{g} \right)$, $\sin \theta_W = g'/\sqrt{g^2 + g'^2}$, and $\cos \theta_W = g/\sqrt{g^2 + g'^2}$ in terms of the couplings g of $SU(2)_L$ and g' of $U(1)_Y$. The weak mixing angle relates the strength of the weak and electromagnetic interaction, which are mediated by the massive W and Z bosons, and the massless photon which is the field quanta of A_μ . Furthermore, the electroweak coupling e of $U(1)_{\text{em}}$ is given by $e \equiv g \sin \theta_W$.

2.4. Electroweak symmetry breaking and the Higgs boson

In the SM, the weak gauge bosons acquire mass via the Higgs mechanism. An arbitrary scalar potential

$$V(\Phi) = \mu^2 \Phi^\dagger \Phi + \lambda (\Phi^\dagger \Phi)^2 \quad (2.14)$$

with a self-interacting SU(2) scalar complex field Φ is introduced, where μ is the mass term of the field. The complex Higgs doublet Φ [1] has four degrees of freedom and it is composed of CP-even and CP-odd components ϕ^0 and a^0 , and the complex charged component ϕ^+ :

$$\Phi = \frac{1}{\sqrt{2}} \begin{pmatrix} \sqrt{2}\phi^+ \\ \phi^0 + ia^0 \end{pmatrix}. \quad (2.15)$$

For $\mu^2 < 0$, the neutral component of the scalar doublet acquires a non-zero vacuum expectation value (VEV) v . Then ϕ^0 can be expanded as $\phi^0 = v + H$ with:

$$\Phi = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + H \end{pmatrix}. \quad (2.16)$$

The $SU(3)_C \times SU(2)_L \times U(1)_Y$ is broken into a $SU(3)_C \times U(1)_{em}$ symmetry. After symmetry breaking, there are three massless Goldstone bosons, which can be identified with three degrees of freedom of the Higgs field. The Higgs field couples to the W_μ and B_μ fields through the covariant derivative in the kinematic term. As a result, the Goldstone bosons mix with the gauge bosons of the corresponding generators of broken symmetries and the physical W and Z gauge bosons acquire masses:

$$M_W^2 = \frac{g^2 v^2}{4}, \quad M_Z^2 = \frac{(g'^2 + g^2) v^2}{4}. \quad (2.17)$$

There is one remaining degree of freedom H, which is the *Higgs boson*. The mass of this new scalar (spin-0) particle is given by $m_H = \sqrt{2\lambda}v$, with the self-coupling parameter λ . The sign of $\mu^2 = -\lambda v^2$ is crucial for the electroweak symmetry to take place. The expectation value of the Higgs field is fixed by the precise measured Fermi coupling G_F : $v = (\sqrt{2}G_F)^{-1/2} \approx 246$ GeV. Also the fermions of the SM acquire masses through the new interactions between the Higgs field and the fermion, the *Yukawa interactions*. In order to allow interactions with up-type quarks, the charge conjugated Higgs doublet $\tilde{\Phi}$ is needed. Then, the Lagrangian of the Yukawa term is given by:

$$\mathcal{L}_{\text{Yukawa}} = -\hat{\lambda}_{d_{ij}} \bar{q}_{L_i} \Phi d_{R_j} - \hat{\lambda}_{u_{ij}} \bar{q}_{L_i} \tilde{\Phi} d_{R_j} - \hat{\lambda}_{\ell_{ij}} \bar{\ell}_{L_i} \Phi e_{R_j} + h.c. , \quad (2.18)$$

where $h.c.$ indicates the Hermitian conjugated terms and $\hat{\lambda}$ are 3×3 matrices for the up and down-type quarks, and the charged leptons. Under electroweak symmetry breaking,

2.4. Electroweak symmetry breaking and the Higgs boson

the Higgs field acquires a VEV and the fermion mass eigenstate basis is chosen, so that the Higgs interactions are diagonalised: $\hat{\lambda}_{f_{ij}} \rightarrow \lambda_{f_i} \delta_{ij}$. The fermions acquire a mass, which is given by $m_{f_i} = \lambda_{f_i} v / \sqrt{2}$, with the corresponding Yukawa coupling λ_f . Furthermore, they couple to the Higgs boson

$$\mathcal{L}_{\text{Yukawa}} = -g_{H_{ff}} \bar{\Psi} \Psi H, \quad (2.19)$$

with $g_{H_{ff}} = \frac{m_f}{v}$. As the masses of the fermions are known, the only free parameter is the mass of the Higgs boson in the SM. The Higgs boson is the most recently discovered fundamental particle. It was discovered by the ATLAS and CMS collaborations in 2012 [3, 4, 13]. Its mass and full width were measured to be $M_H = 125.09 \pm 0.24$ GeV and $\Gamma < 1.7$ GeV [1].

2.5. Main shortcomings of the Standard Model

The Standard Model is a very successful theory and the electroweak theory has been verified over a wide range of energies. It was crowned with the discovery of the Higgs boson by the ATLAS and CMS collaborations in 2012 [3, 4, 13]. However, the SM is an incomplete theory as it does not incorporate gravitation.

In the original formulation of the SM, neutrinos are strictly massless, because there is no Dirac mass term due to the absence of right-handed neutrinos and there is no Majorana mass term due to conservation of baryon and lepton number. However, neutrino experiments have established the existence of neutrino oscillations, transitions between the different neutrino flavours. This is an evidence for neutrino mixing and nonzero neutrino masses [1].

Astrophysical observations imply the existence of an unknown type of matter and energy, called dark matter and dark energy, which cannot be explained within the SM. Space observatories like Planck have studied the fluctuations in the spectrum of the relic cosmic microwave background (CMB) and have established the existence of cold dark matter [14]. In 2011, the Nobel Prize in Physics was awarded to S. Perlmutter, B.P. Schmidt, and A.G. Riess for "the discovery of the accelerating expansion of the Universe through observations of distant supernovae" [15]. In addition, the CP-violating effects of the SM cannot explain the large difference between matter and anti-matter in the universe [1].

Furthermore, the flavour structure is not explained in the theoretical description of the SM [2]. There is a large number of free parameters, such as masses and mixing parameters. The Higgs mechanism can explain how masses are introduced in the SM, but it cannot explain its values.

3. Lepton-flavour violating decays of the Higgs boson

The Standard Model it is considered to be an effective theory up to a certain scale Λ . One shortcoming is that the flavour structure is not incorporated in the theoretical description of the SM. There are three main puzzles in the SM related to flavour. First, the replication of the generation constitutes. Secondly, there is no principle underlying the form of the Yukawa matrices describing the SM Yukawa interactions. Thirdly, the peculiar pattern of fermion masses and mixing cannot be explained [2]. New physics could explain some of the symmetries and open questions [2]. Grand Unification theories (GUTs) try to unify all known interactions but gravity in a single gauge group, while extra dimensions could link flavour to the geometry of this extra dimensions. A simple extension of the Higgs sector with at least two Higgs doublets could add additional CP-violation.

A Higgs boson was discovered in 2012, whose properties are compatible with the SM Higgs boson, within experimental uncertainties. Its mass has been measured by the ATLAS and CMS collaborations and found to be 125.09 ± 0.24 GeV [1]. After the discovery of the Higgs boson, its nature has to be investigated by measuring its properties including its couplings to charged leptons. Charged leptons acquire mass through Yukawa interactions with the Higgs field, as described in Chapter 2. In the SM, the Yukawa interaction matrix can be diagonalised in the mass basis. Thus, lepton-flavour violating (LFV) decays of the Higgs boson are forbidden. Extensions of the SM, such as multi-Higgs doublet models, can introduce LFV Yukawa couplings. Multi-Higgs doublet models have in general tree level Higgs-mediated flavour changing neutral currents leading to lepton-flavour violating (LFV) Yukawa couplings [2]. LFV Yukawa couplings can also be introduced in composite Higgs models [16] and models with extra dimensions [6].

In the following, a short review on lepton-flavour violating decays of the Higgs bosons is given, which was the status at time when the search was started. It is described how lepton-flavour violating Yukawa-couplings can be introduced by new physics and how they are constrained by low-energy measurements [6]. Finally, it is described how LFV decays of the Higgs boson can be measured at the LHC and which sensitivity is expected, including observed upper limits on the branching fraction [17].

3.1. Introduction of lepton-flavour violating decays

The Higgs mechanism is responsible for the masses of the gauge bosons and the fermions. First, it is discussed how charged leptons gain mass in the SM. Interactions between charged leptons with the Higgs field are mediated by Yukawa interactions, therefore the structure of the corresponding Yukawa matrix in the mass basis is afterwards discussed. Then, it is explained how LFV Yukawa couplings can be introduced considering an effective field theory approach [6]. Finally, an overview is given on the structure of the Yukawa matrix for charged leptons in general. Decay widths and branching fractions of the LFV decays in terms of the corresponding Yukawa couplings are also shown.

Yukawa couplings of charged leptons in the SM Charged leptons acquire a mass via the Higgs mechanism, which was introduced in Section 2.4. A direct mass term is not gauge invariant, but interactions of the charged leptons ℓ to the scalar Higgs field ϕ can be introduced with the following term:

$$\mathcal{L}_{\text{SM}} = -\lambda_{ij} \bar{L}_L^i \Phi \ell_R^j + h.c. , \quad (3.1)$$

where L_L^i , with $\bar{L}_L^i = (\bar{\nu}_\ell^i, \bar{\ell}_L^i)$, are the $\text{SU}(2)_L$ doublets, and ℓ_R^i the weak singlets with the indices i, j running over generations.

After electroweak symmetry breaking, the Higgs doublet can be written in terms of the vacuum expectation value v and the physical Higgs boson H :

$$\Phi = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + H \end{pmatrix}, \quad (3.2)$$

leading to the following Lagrangian:

$$\begin{aligned} \mathcal{L}_{\text{SM}} &= -\frac{\lambda_{ij}}{\sqrt{2}} (v + H) \bar{\ell}_L^i \ell_R^j + h.c. \\ &= -\frac{\lambda_{ij}}{\sqrt{2}} v \boxed{\bar{\ell}_L^i \ell_R^j} - \frac{\lambda_{ij}}{\sqrt{2}} \boxed{\bar{\ell}_L^i \ell_R^j H} + h.c., \end{aligned} \quad (3.3)$$

where the term proportional to $\bar{\ell}_L^i \ell_R^j$ is the mass term of the charged leptons and the term proportional to $\bar{\ell}_L^i \ell_R^j H$ describes the Yukawa interactions. The mass matrix m , with $m_{ij} = \frac{\lambda_{ij}}{\sqrt{2}} v$, can be diagonalised using the matrices V_L and V_R :

$$m = \begin{pmatrix} m_e & 0 & 0 \\ 0 & m_\mu & 0 \\ 0 & 0 & m_\tau \end{pmatrix} = \frac{1}{\sqrt{2}} V_L \lambda V_R^\dagger v = \frac{\lambda^m}{\sqrt{2}} v, \quad (3.4)$$

where $\lambda^m = V_L \lambda V_R^\dagger$ is written in terms of the mass basis. The Yukawa interaction matrix Y can also be written in terms of the mass basis:

$$Y = V_L \frac{\lambda}{\sqrt{2}} V_R^\dagger = \frac{\lambda^m}{\sqrt{2}} = \frac{m}{v} = \begin{pmatrix} Y_{ee} & 0 & 0 \\ 0 & Y_{\mu\mu} & 0 \\ 0 & 0 & Y_{\tau\tau} \end{pmatrix}, \quad (3.5)$$

which is also diagonal. Thus, there are no lepton-flavour violating (LFV) Yukawa interactions in the SM.

Dimension-6 operators If the Higgs field is the only field that causes electroweak symmetry breaking and if the particle spectrum is consisting only out of SM particles, including the Higgs boson, up to some energy scale $\Lambda \gg 200$ GeV, additional heavy fields can be integrated out, leading to an effective field theory.

Terms including higher-dimension operators can be added to the SM Lagrangian, which has dimension four. An effective Lagrangian can be written including dimension-6 operators O_n^{ij} :

$$\mathcal{L}_{eff} = \mathcal{L}_{SM} + \sum_{nij} \frac{\alpha_n^{ij}}{\Lambda^2} O_n^{ij}, \quad (3.6)$$

where i, j denote flavour indices, n runs over the number of independent operators, Λ is the scale of new physics, and α_n^{ij} are coefficients [18]. This effective Lagrangian describes the effects of new physics at the electroweak scale in a model-independent way.

Neutrino masses can be generated by a dimension-5 operator [18]. However, this operator does not generate LFV interactions for the Higgs boson.

The dimension-6 operator $O_{L\phi}^{ij} = (\Phi^\dagger \Phi)(\bar{L}_i \ell_{Rj} \Phi)$ is a Yukawa-type operator that only generates LFV interactions of the Higgs boson. It also contributes to the fermion mass matrices [18]. These contributions have to be included in the fermion mass matrices. Then, the mass matrices are diagonalised. Finally, the interaction term of the Lagrangian is written in terms of the mass basis. Including the Yukawa-type operator to the Lagrangian leads to the following additional term:

$$\Delta \mathcal{L} = -\frac{\lambda'_{ij}}{\Lambda^2} (\Phi^\dagger \Phi)(\bar{L}_i^j \ell_{Rj}^i \Phi) + h.c., \quad (3.7)$$

where λ'_{ij} is the coefficient α_n^{ij} of the dimension-6 operator $O_{L\phi}^{ij}$.

3. Lepton-flavour violating decays of the Higgs boson

Electroweak symmetry breaking leads to the following Lagrangian terms:

$$\begin{aligned}
\Delta\mathcal{L} &= -\frac{\lambda'_{ij}}{\Lambda^2} \left(\frac{v+H}{\sqrt{2}} \right)^2 \bar{\ell}_L^i \ell_R^j \left(\frac{v+H}{\sqrt{2}} \right) \\
&= -\frac{\lambda'_{ij}}{\Lambda^2} \frac{1}{(\sqrt{2})^3} (v^3 + 3v^2H + 3vH^2 + H^3) \bar{\ell}_L^i \ell_R^j \\
&= -\frac{1}{\sqrt{2}} \frac{\lambda'_{ij}}{2\Lambda^2} \left(v^3 \boxed{\bar{\ell}_L^i \ell_R^j} + 3v^2 \boxed{\bar{\ell}_L^i \ell_R^j H} + 3v \bar{\ell}_L^i \ell_R^j H^2 + \bar{\ell}_L^i \ell_R^j H^3 \right),
\end{aligned} \tag{3.8}$$

where term proportional to $\bar{\ell}_L^i \ell_R^j$ leads to an additional mass term of the charged leptons and the term proportional to $\bar{\ell}_L^i \ell_R^j H$ adds a further Yukawa-interaction term. Thus, the mass term of the effective Lagrangian $\mathcal{L}_{\text{mass}}$ is the sum of the SM mass term and the additional contribution:

$$\begin{aligned}
\mathcal{L}_{\text{mass}} &= \mathcal{L}_{\text{mass,SM}} + \Delta\mathcal{L}_{\text{mass}} \\
&= -\left(\frac{\lambda_{ij}}{\sqrt{2}} v + \frac{1}{\sqrt{2}} \frac{\lambda'_{ij}}{2\Lambda^2} v^3 \right) \bar{\ell}_L^i \ell_R^j \\
&= -\frac{1}{\sqrt{2}} \left(\lambda_{ij} + \frac{v^2}{2\Lambda^2} \lambda'_{ij} \right) v \bar{\ell}_L^i \ell_R^j.
\end{aligned} \tag{3.9}$$

The term of the effective Lagrangian which describes the Yukawa interactions, $\mathcal{L}_Y$, is given by:

$$\begin{aligned}
\mathcal{L}_Y &= \mathcal{L}_{Y,\text{SM}} + \Delta\mathcal{L}_Y \\
&= -\left(\frac{\lambda_{ij}}{\sqrt{2}} + \frac{1}{\sqrt{2}} \frac{\lambda'_{ij}}{2\Lambda^2} 3v^2 \right) \bar{\ell}_L^i \ell_R^j H \\
&= -\frac{1}{\sqrt{2}} \left(\lambda_{ij} + 3 \frac{v^2}{2\Lambda^2} \lambda'_{ij} \right) \bar{\ell}_L^i \ell_R^j H.
\end{aligned} \tag{3.10}$$

The mass matrix of the charged leptons can be diagonalised as for the SM using other matrices V_L, V_R :

$$\sqrt{2}m = V_L \left[\lambda + \frac{v^2}{2\Lambda^2} \lambda' \right] V_R^\dagger v. \tag{3.11}$$

Then, the Yukawa interaction matrix can be written in terms of the mass basis:

$$\begin{aligned}
\sqrt{2}Y &= V_L \left[\lambda + 3 \frac{v^2}{2\Lambda^2} \lambda' \right] V_R^\dagger \\
&= \frac{\sqrt{2}m}{v} + V_L 2 \frac{v^2}{2\Lambda^2} \lambda' V_R^\dagger \\
&= \frac{\sqrt{2}m}{v} + \frac{v^2}{\Lambda^2} V_L \lambda' V_R^\dagger \\
&= \frac{\sqrt{2}m}{v} + \frac{v^2}{\Lambda^2} \hat{\lambda},
\end{aligned} \tag{3.12}$$

3.1. Introduction of lepton-flavour violating decays

where m is the diagonal mass matrix and $\hat{\lambda} = V_L \lambda' V_R^\dagger$. The first term is the Yukawa term of the SM, while the second term can introduce non-diagonal terms as $\hat{\lambda}$ is in principle an arbitrary non-diagonal matrix [6], leading to Yukawa couplings Y_{ij} given by:

$$Y_{ij} = \frac{m_i}{v} \delta_{ij} + \frac{v^2}{\sqrt{2}\Lambda^2} \hat{\lambda}_{ij}, \quad (3.13)$$

with possible LFV Yukawa interactions.

Lepton-flavour violating Yukawa couplings The inclusion of LFV Yukawa interactions was discussed by introducing a dimension-6 operator to the Lagrangian. There are several possible additional terms to the Lagrangian due to new physics which prevent the diagonalisation of the mass matrix and the Yukawa matrix at the same time. Another example are theories with more than one Higgs doublet, where the scalar fields can mix leading to a neutral Higgs boson with in general LFV Yukawa interactions. Thus, the Yukawa term of the Lagrangian for the charged leptons in the mass basis is after the electroweak symmetry breaking in general given by:

$$\mathcal{L}_{\text{Yukawa}} = -m_i \bar{\ell}_L^i f_R^i - Y_{ij} \bar{\ell}_L^i \ell_R^j H + h.c. , \quad (3.14)$$

where Y_{ij} are the entries of the Yukawa matrix:

$$Y = \begin{pmatrix} Y_{ee} & Y_{e\mu} & Y_{e\tau} \\ Y_{\mu e} & Y_{\mu\mu} & Y_{\mu\tau} \\ Y_{\tau e} & Y_{\tau\mu} & Y_{\tau\tau} \end{pmatrix}, \quad (3.15)$$

which can have LFV Yukawa couplings Y_{ij} . LFV Yukawa couplings allow lepton-flavour violating decays of the Higgs boson. The decay width of LFV Higgs-boson decays is

$$\Gamma(H \rightarrow \ell_i \ell_j) = \frac{m_H}{8\pi} (|Y_{ji}|^2 + |Y_{ij}|^2). \quad (3.16)$$

Using the decay width of the LFV decay, the branching fraction of the LFV Higgs-boson decay can be derived, when contributions of the other LFV decays are negligible

$$B(H \rightarrow \ell_i \ell_j) = \frac{\Gamma(H \rightarrow \ell_i \ell_j)}{\Gamma(H \rightarrow \ell_i \ell_j) + \Gamma_{\text{SM}}}, \quad (3.17)$$

where Γ_{SM} is the total width of the Higgs boson in the SM. Non-negligible contributions to the total decay width from other LFV Higgs boson decays would have to be added to Γ_{SM} for estimating the branching fractions.

3. Lepton-flavour violating decays of the Higgs boson

In the SM, the Higgs boson couples to all particles according to their mass. In presence of new physics, LFV Yukawa couplings might be related closely to the fermion mass matrices, reflecting the observed fermion mass hierarchy [19]. Then, the LFV couplings have a hierarchical structure given by $\Delta_{ij}\sqrt{m_i m_j}$, where i, j are the generation indices, and Δ_{ij} is related to the mass mixing. Mass corrections to the lepton masses coming from mass mixing are assumed to be small in the so called *natural* assumption, where SM Yukawa couplings are taken for the flavour-diagonal couplings [16]. Thus, LFV branching fractions can be expressed in terms of SM ones:

$$B_{nat}(H \rightarrow ij) = B_{SM}(H \rightarrow ii) \cdot \frac{m_j}{m_i}, \quad (3.18)$$

with $\Gamma_{tot} \approx \Gamma_{SM}$, $B_{nat}(H \rightarrow ij)$ being the *natural branching fraction* and $B_{SM}(H \rightarrow ii)$ being the branching fraction for the SM Higgs-boson decay to two leptons i . The LFV branching fractions are expected to be of same size or smaller than $B_{nat}(H \rightarrow ij)$. Larger total widths lead to a smaller branching fraction. A naturalness assumption [6], referred to as *theoretical naturalness limit*, can also be derived for the Yukawa couplings Y_{ij} [6]:

$$|Y_{ji}Y_{ij}| \lesssim \frac{m_i m_j}{v^2}, \quad (3.19)$$

with v being the vacuum expectation value of the Higgs boson.

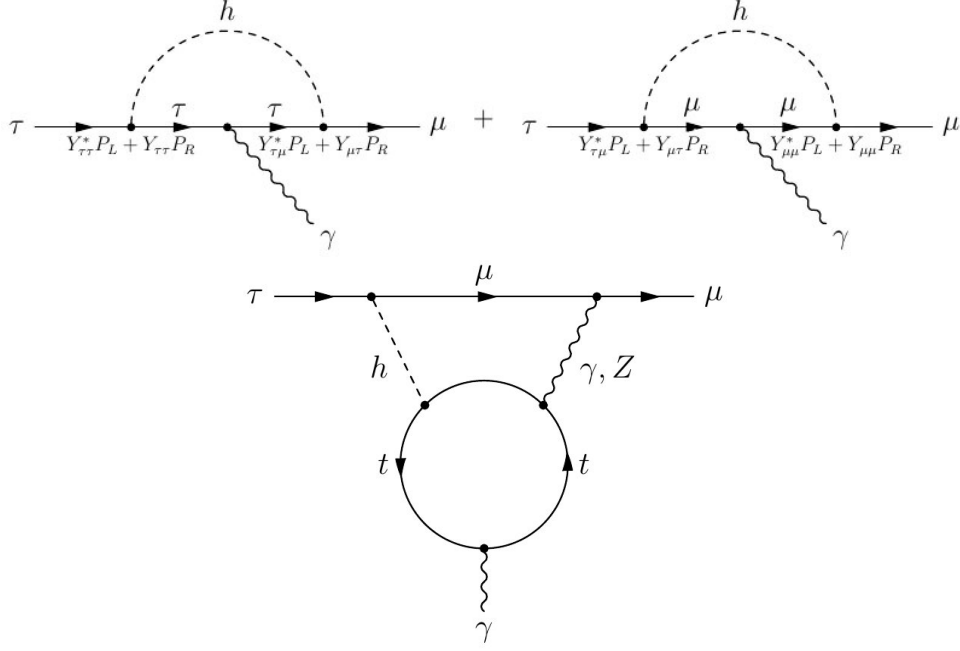


Figure 3.1.: Diagrams contributing to the lepton-flavour violating decay $\tau \rightarrow \mu\gamma$, mediated by a Higgs boson [6]. Similar for $\tau \rightarrow e\gamma$.

3.2. Constraints from low-energy measurements

Constraints on lepton-flavour violating decays of the Higgs boson have been derived in [6] using low-energy measurements and certain assumptions on the value of flavour-diagonal Yukawa couplings. In the following, constraints from several low-energy measurements on LFV Yukawa-couplings are discussed. Diagrams which can contribute to the low-energy measurements are shown. The results of the measurements and assumptions used to derive constraints on the LFV Yukawa-couplings [6] are given. Finally, an overview of all constraints from [6] is given.

Constraints from LFV decays $\ell_i \rightarrow \ell_j\gamma$ Lepton-flavour violating Yukawa couplings contribute also to LFV decays of type $\ell_i \rightarrow \ell_j\gamma$, where i, j are flavour indices with $i \neq j$. One-loop diagrams and two-loop diagrams of this process are shown in Figure 3.1. Constraints on the LFV Yukawa couplings of the Higgs boson with $m_H = 125$ GeV can be derived assuming SM values for $Y_{\tau\tau}$, $Y_{\mu\mu}$, and Y_{tt} . Experimental upper bounds on the branching fractions $B(\ell_i \rightarrow \ell_j\gamma)$ used for deriving constraints are summarised in Table 3.1. Constraints of type $\sqrt{|Y_{ij}|^2 + |Y_{ji}|^2}$ can be derived. A further constraint on $(|Y_{\tau\mu}Y_{e\tau}|^2 + Y_{\mu\tau}Y_{\tau e}|^2)^{1/4}$ can be obtained from $\mu \rightarrow e\gamma$ by setting $Y_{\mu e}$ and $Y_{e\mu}$ to zero.

3. Lepton-flavour violating decays of the Higgs boson

Table 3.1.: Experimental upper bounds on the branching fractions of LFV decays $\ell_i \rightarrow \ell_j \gamma$ used in [6], originally taken from [20, 21].

Decay	$B(\tau \rightarrow \mu \gamma)$	$B(\tau \rightarrow e \gamma)$	$B(\mu \rightarrow e \gamma)$
Experimental bound	$< 4.4 \times 10^{-8}$	$< 3.3 \times 10^{-8}$	$< 2.4 \times 10^{-12}$

Constraints from LFV decays $\ell_i \rightarrow 3\ell_j$ LFV decays of type $\ell_i \rightarrow 3\ell_j$ can constrain the LFV Yukawa couplings. Figure 3.2 shows the tree-level contribution for a Higgs boson exchange and higher-order contributions, similar to decays of type $\ell_i \rightarrow \ell_j \gamma$ with the photon decaying to two leptons. Such decays have been searched for and the experimental upper bounds on the branching fractions of are summarised in Table 3.2. The previous assumptions on the mass of the Higgs boson and flavour-diagonal Yukawa couplings were used to derive constraints on the LFV Yukawa couplings. Z boson contribution to the LFV decay $\tau \rightarrow \mu$, described in [22] are neglected, as it turns out that their contributions are negligible [6, 22].

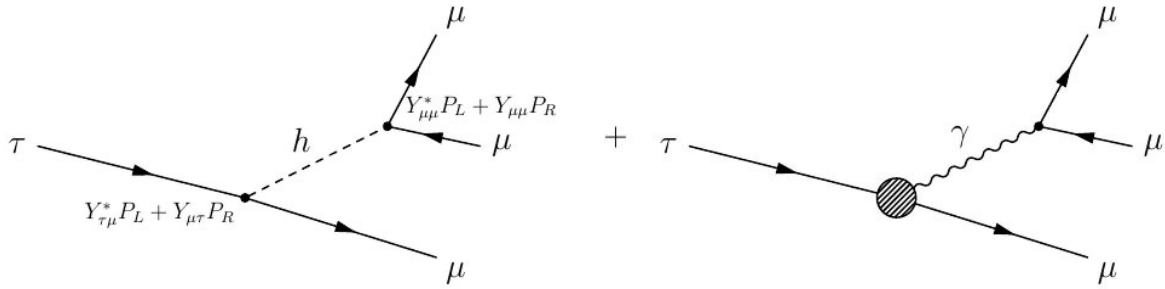


Figure 3.2.: Diagrams contributing to the LFV decay $\tau \rightarrow 3\mu$ [6]. The tree-level Higgs boson exchange contribution (left) and the higher-order contribution (right), where the blob represents loops of the form shown in Figure 3.1. Similar for $\tau \rightarrow 3e$.

Table 3.2.: Experimental upper bounds on the branching fractions of LFV decays $\ell_i \rightarrow 3\ell_j$ used in [6], originally taken from [23, 24].

Decay	$B(\tau \rightarrow 3\mu)$	$B(\tau \rightarrow 3e)$	$B(\mu \rightarrow 3e)$
Experimental bound	$< 2.1 \times 10^{-8}$	$< 2.7 \times 10^{-8}$	$< 1.0 \times 10^{-12}$

Constraints from muonium-antimuonium oscillations The bound state $\mu^+ e^-$ is called *muonium* (M) and can oscillate to the $e^+ \mu^-$ bound state called antimuonium ($\bar{M}$), shown in Figure 3.3. An upper constraint on the conversion probability $P(M \rightarrow \bar{M}) < 8.3 \times 10^{-11}$ [25] was derived by the muonium-antimuonium conversion spectrometer (MACS) experiment at PSI. The time-integrated conversion probability depends on the mass splitting between the two mass eigenstates of the mixed $M - \bar{M}$ system, which depends on

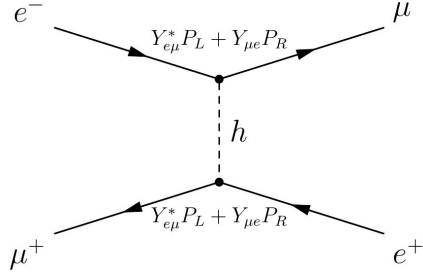


Figure 3.3.: Diagram leading to muonium-antimuonium oscillations [6].

$|Y_{\mu e} + Y_{e\mu}^*|$ [6]. A correction factor $S_B = 0.35$, accounting for the splitting of the muonium states in the magnetic field of the detector, was applied on the conversion probability $P(M \rightarrow \bar{M}) < 8.3 \times 10^{-11}/S_B$ when deriving the constraints on the LFV Yukawa couplings [6], leading to a weaker upper bound on the conversion probability. The correction factor depends on the conversion operator and the smallest value was used in [6] to derive the constraints.

Constraints from magnetic and electric dipole moments The experimental value of the magnetic dipole moment of the muon $g_\mu - 2$ is more than 3 standard deviations above the SM prediction. LFV Yukawa couplings can contribute to the magnetic moment via diagrams shown in Figure 3.4, which also contribute to the electric dipole moment. Neglecting terms suppressed by m_μ/m_τ or m_τ/m_H , the LFV contribution to $g_\mu - 2$ due to the one-loop diagram has the form [6]

$$a_\mu \equiv \frac{g_\mu - 2}{2} \propto \text{Re}(Y_{\mu\tau} Y_{\tau\mu}), \quad (3.20)$$

with the discrepancy between measurement and SM prediction [6] having the size

$$\Delta a_\mu \equiv a_\mu^{\text{exp}} - a_\mu^{\text{SM}} = (2.87 \pm 0.63 \pm 0.49) \times 10^{-9}, \quad (3.21)$$

where a_μ^{exp} and a_μ^{SM} are the measured and predicted value, respectively¹. The one-loop diagram in Figure 3.4 can contribute to an electric dipole moment (EDM) of the muon if the LFV Yukawa couplings are complex. If the terms suppressed by m_μ/m_τ or m_τ/m_H are neglected, the dependence on the LFV Yukawa couplings of the electric dipole moment d_μ is

$$d_\mu \propto -\text{Im}(Y_{\mu\tau} Y_{\tau\mu}). \quad (3.22)$$

¹The updated value from the Particle Data is [1]: $\Delta a_\mu = (2.88 \pm 0.63 \pm 0.49) \times 10^{-9}$

3. Lepton-flavour violating decays of the Higgs boson

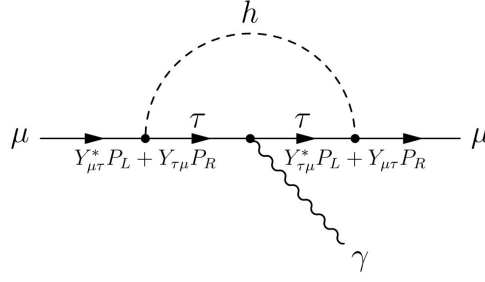


Figure 3.4.: Diagram contributing to magnetic and electric dipole moments of the muon [6].

Constraints from $\mu \rightarrow e$ conversion in nuclei LFV Yukawa couplings can contribute to $\mu \rightarrow e$ conversion in nuclei via tree-level exchange of a Higgs boson and one-loop diagrams with a Higgs boson and a photon exchange. These diagrams are shown in Figure 3.5. Two-loop contributions are larger than the one-loop ones, as they are only suppressed by the weak gauge coupling or Y_{tt} . Thus, they were taken into account for deriving the constraints on the LFV Yukawa couplings [6] assuming SM values for the flavour-diagonal Yukawa couplings, as described before. These Yukawa couplings have not been observed at the LHC up to today. Using these assumption, the LFV decays of type $\ell_i \rightarrow \ell_j \gamma$ give the strongest constraints.

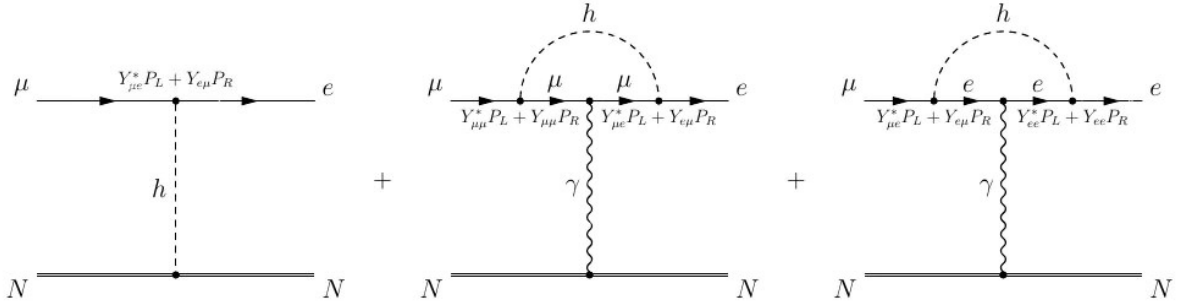


Figure 3.5.: Diagrams contributing to $\mu \rightarrow e$ conversion in nuclei [6]. The tree-level diagram (left) with an exchange of an Higgs boson and one-loop diagrams (middle and right) with a Higgs-boson loop and a photon exchange can contribute.

Constraints on the LFV Yukawa couplings from low-energy measurements are summarised in Table 3.3 and Figure 3.6. They were in most cases derived [6] using certain assumptions regarding the flavour-diagonal Yukawa couplings. These measurements constrain the branching fractions of LFV Higgs-boson decays to $\mu\tau$ or $e\tau$ to be $\lesssim \mathcal{O}(10^{-1})$, while the constraint for the decay to μe is stronger and $\lesssim \mathcal{O}(10^{-8})$. Constraints on the branching fractions were derived under the assumption that only one of them contributes to the total width of the Higgs boson additionally to the SM total width.

Table 3.3.: Constraints on lepton-flavour violating Yukawa couplings from low-energy measurements. The constraint from $M - \bar{M}$ oscillations was derived applying a correction factor $1/S_B = 1/0.35$ on the conversion probability. The mass of the observed Higgs boson, $m_H = 125$ GeV, was used to derive constraints from the other low-energy measurements. Mass terms suppressed by m_μ/m_τ or m_τ/m_H were neglected for deriving constraints from the measurements of the magnetic or electric dipole momentum of the muon and electron, respectively. SM values were assumed for the flavour-diagonal Yukawa couplings $Y_{\tau\tau}$, $Y_{\mu\mu}$, and Y_{tt} when deriving constraints from measurements of type $\ell_i \rightarrow \ell_j\gamma$, $\ell_i \rightarrow 3\ell_j$, and $\mu \rightarrow e$ conversion in nuclei. The constraint on $(|Y_{\tau\mu}Y_{e\tau}|^2 + |Y_{\mu\tau}Y_{\tau e}|^2)^{1/4}$ was derived setting $Y_{\mu e}$ and $Y_{e\mu}$ to zero. Instead of an upper bound on the Yukawa couplings from the magnetic dipole measurement of the muon, the needed value is listed for explaining the deviation from the SM prediction. These constraints on the LFV Yukawa couplings and how they are derived were presented in [6].

Channel	Coupling	Bound
$\mu \rightarrow e\gamma$	$\sqrt{ Y_{\mu e} ^2 + Y_{e\mu} ^2}$	$< 3.6 \times 10^{-6}$
$\mu \rightarrow 3e$	$\sqrt{ Y_{\mu e} ^2 + Y_{e\mu} ^2}$	$\lesssim 3.1 \times 10^{-5}$
electron $g - 2$	$\text{Re}(Y_{e\mu}Y_{\mu e})$	$-0.019...0.026$
electron EDM	$ \text{Im}(Y_{e\mu}Y_{\mu e}) $	$< 9.8 \times 10^{-8}$
$\mu \rightarrow e$ conversion	$\sqrt{ Y_{\mu e} ^2 + Y_{e\mu} ^2}$	$< 1.2 \times 10^{-5}$
$M - \bar{M}$ oscillations	$ Y_{\mu e} + Y_{e\mu}^* $	< 0.079
$\tau \rightarrow e\gamma$	$\sqrt{ Y_{\tau e} ^2 + Y_{e\tau} ^2}$	< 0.014
$\tau \rightarrow 3e$	$\sqrt{ Y_{\tau e} ^2 + Y_{e\tau} ^2}$	$\lesssim 0.12$
electron $g - 2$	$\text{Re}(Y_{e\tau}Y_{\tau e})$	$[-2.1...2.9] \times 10^{-3}$
electron EDM	$ \text{Im}(Y_{e\tau}Y_{\tau e}) $	$< 1.1 \times 10^{-8}$
$\tau \rightarrow \mu\gamma$	$\sqrt{ Y_{\tau\mu} ^2 + Y_{\mu\tau} ^2}$	< 0.016
$\tau \rightarrow 3\mu$	$\sqrt{ Y_{\tau\mu} ^2 + Y_{\mu\tau} ^2}$	$\lesssim 0.25$
muon $g - 2$	$\text{Re}(Y_{\mu\tau}Y_{\tau\mu})$	$(2.7 \pm 0.75) \times 10^{-3}$
muon EDM	$ \text{Im}(Y_{\mu\tau}Y_{\tau\mu}) $	$-0.8...1.0$
$\mu \rightarrow e\gamma$	$(Y_{\tau\mu}Y_{e\tau} ^2 + Y_{\mu\tau}Y_{\tau e} ^2)^{1/4}$	$< 3.4 \times 10^{-4}$

3. Lepton-flavour violating decays of the Higgs boson

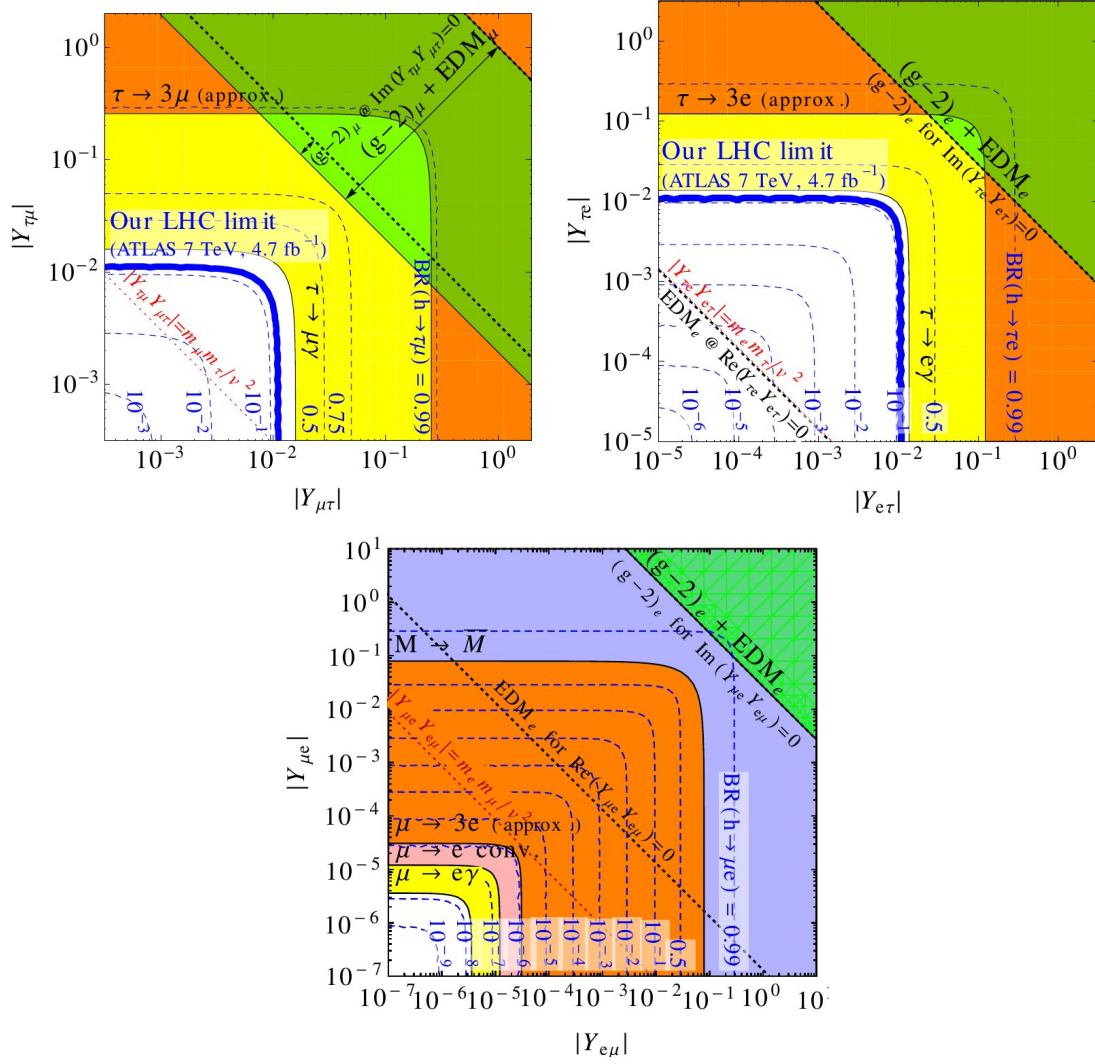


Figure 3.6.: Constraints on the lepton-flavour violating Yukawa couplings for a 125 GeV Higgs boson [6]. The diagonal Yukawa couplings are approximated by their SM values. Thin blue dashed lines are contours of constant branching fraction, whereas thick blue lines are LHC constraints derived in [6]. Shaded regions show the constraints from low-energy measurements. The thin red dotted lines show the theoretical naturalness limits $|Y_{ji}Y_{ij}| \lesssim m_i m_j / v^2$.

3.3. Lepton-flavour violating decays of the Higgs boson at the LHC

A Higgs boson was discovered in 2012, whose properties are compatible with the SM Higgs boson, within experimental uncertainties. Its mass has been measured by the ATLAS and CMS collaborations and found to be 125.09 ± 0.24 GeV [1]. The nature of the discovered Higgs boson has to be investigated by measuring its properties including its couplings to charged leptons.

In the following, SM production-mechanisms of the Higgs boson are discussed, which are relevant at the LHC. Possible lepton-flavour violating decays of the Higgs boson are presented. The size of the branching fractions on SM Higgs-boson decays and the natural LFV branching fractions, Section 3.1, is discussed and a status is given about the upper limits on the branching fraction of these decays, which was the status in July 2014.

Higgs-boson production mechanisms Figure 3.7 shows cross sections of different Higgs-boson production-mechanisms in the SM for proton-proton collisions and a Higgs boson with a mass of 125 GeV. Cross sections are shown as function of centre-of-mass energy $\sqrt{s}$ in a range that is relevant at the LHC. Gluon fusion (GF) and vector-boson fusion (VBF) are the dominant Higgs-boson production-mechanisms. Other production mechanisms have a lower cross section, therefore they can be neglected for this search.

Lepton-flavour violating Yukawa couplings Lepton-flavour violating Yukawa couplings can introduce decays of the Higgs boson to $\mu\tau$, $e\tau$, or $e\mu$. Figure 3.8 shows the Feynman diagrams for lepton-flavour violating decays of the Higgs boson. Decays of the Higgs boson to $\mu\tau$ or $e\tau$ can be split into a leptonic channel (τ_e/τ_μ) and a hadronic channel (τ_h). The branching fractions of a tau to a muon or an electron are $\sim 20\%$, while $\sim 65\%$ of the taus decay hadronically [1].

Taus have a short life-time of $\tau = (290.3 \pm 0.5) \times 10^{-15}$ s and a decay length of $c\tau = 87.03$ μm [1] leading to a secondary vertex, which can be resolved within the CMS pixel tracking-detector. Muons, electrons, and hadronic taus can be reconstructed with the CMS detector. Neutrinos from tau decays lead to missing transverse energy $\cancel{E}_T$.

Feynman diagrams for $H \rightarrow \mu\tau_e$ (a), $H \rightarrow e\tau_\mu$ (c), and $H \rightarrow e\mu$ (e) lead to a signature with an isolated electron and an isolated muon of opposite charge in the CMS detector. The last decay mode can be distinguished from the other Higgs-boson decays by absence of $\cancel{E}_T$ due to neutrinos. Decays to $\mu\tau_e$ or $e\tau_\mu$ can be distinguished by the small angle between the direction of the lepton from the tau decay and the direction of the missing transverse momentum $\vec{\cancel{E}}_T$, due to the neutrinos, for high momenta of the tau. Feynman diagrams for $H \rightarrow \mu\tau_h$ (b), or $H \rightarrow e\tau_h$ (d) have signatures with a hadronic tau and an isolated

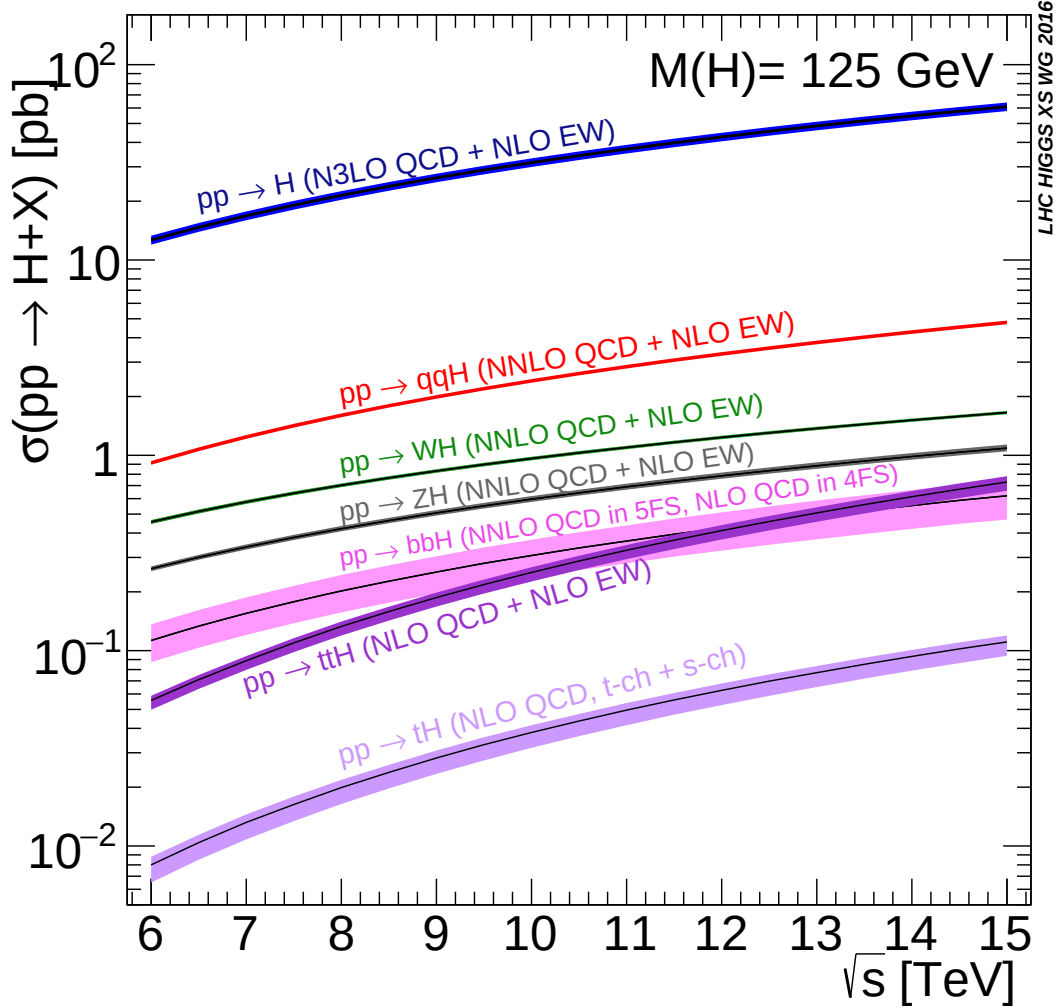


Figure 3.7.: Production cross sections for a Higgs boson with a mass of 125 GeV. The production cross sections have been calculated for proton-proton collisions as function of the centre-of mass energy $\sqrt{s}$ in a range that is relevant for the LHC. Gluon-fusion production (blue line) and vector-boson fusion (red line) are the dominant Higgs-boson production mechanisms. Cross sections for associated Higgs-boson production, WH (dark-green line), ZH (black line), $b\bar{b}H$ Higgs-boson production (pink line), and $t\bar{t}H$ or tH production (violet lines) are several orders of magnitude smaller than both dominant production mechanisms [26, 27].

3.3. Lepton-flavour violating decays of the Higgs boson at the LHC

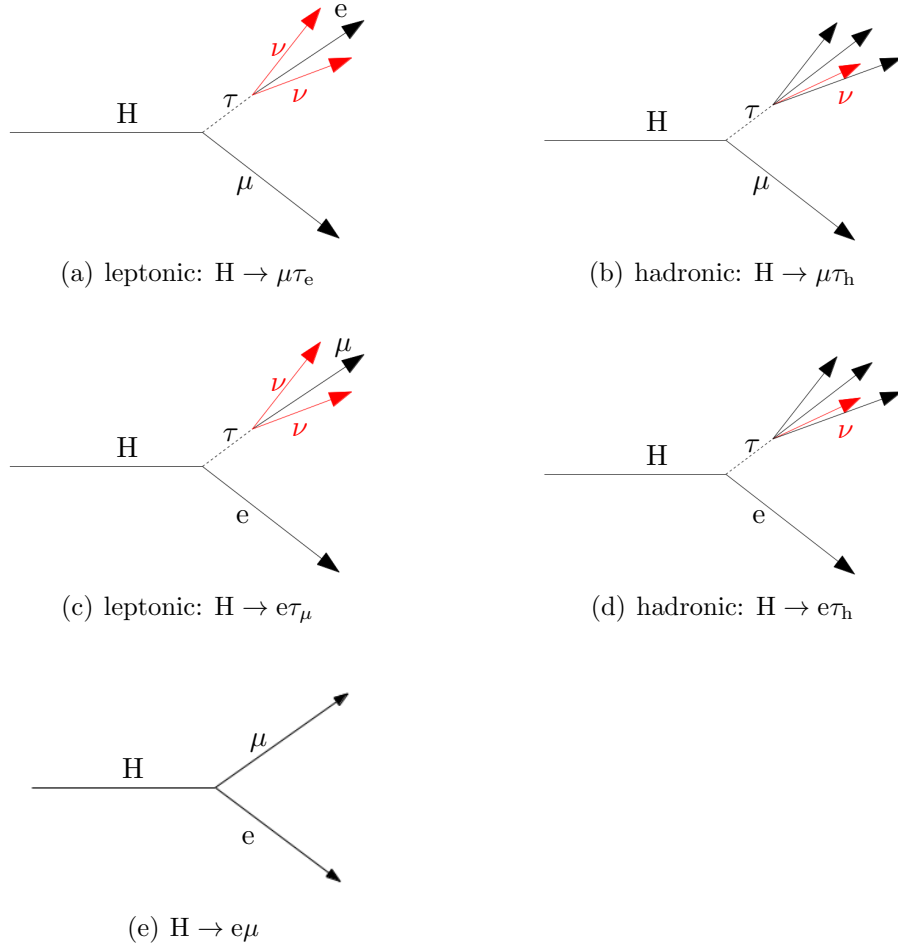


Figure 3.8.: Feynman diagrams of lepton-flavour violating Higgs-boson decays. First row shows diagrams for the coupling of the Higgs boson to $\mu\tau$ (a,b). Couplings to $e\tau$ (c,d) and to $e\mu$ (e) are shown in the second and third row, respectively. Taus can decay leptonically or hadronically. Feynman diagrams are shown for the leptonic decay of taus (a,c) and for the hadronic decay of taus (b,d).

electron or muon, respectively. Their signature include $\cancel{E}_T$ as well due to the neutrino of the tau decay. Tau decays to a lepton which has the same flavour as the lepton of the Higgs-boson decay, the channels $e\tau_e$ and $\mu\tau_\mu$, are not considered because their signatures are very similar to Z-boson decays leading to an overwhelming background of di-leptons for such processes.

Branching fractions In the SM, the Higgs boson should couples to all particles according to their mass. Figure 3.9 shows the branching fractions for the SM Higgs-boson. The branching fraction for Higgs-boson decays to taus is in the order of 6%, and the branching fraction of the Higgs boson to muons is in the order of 0.02%.

3. Lepton-flavour violating decays of the Higgs boson

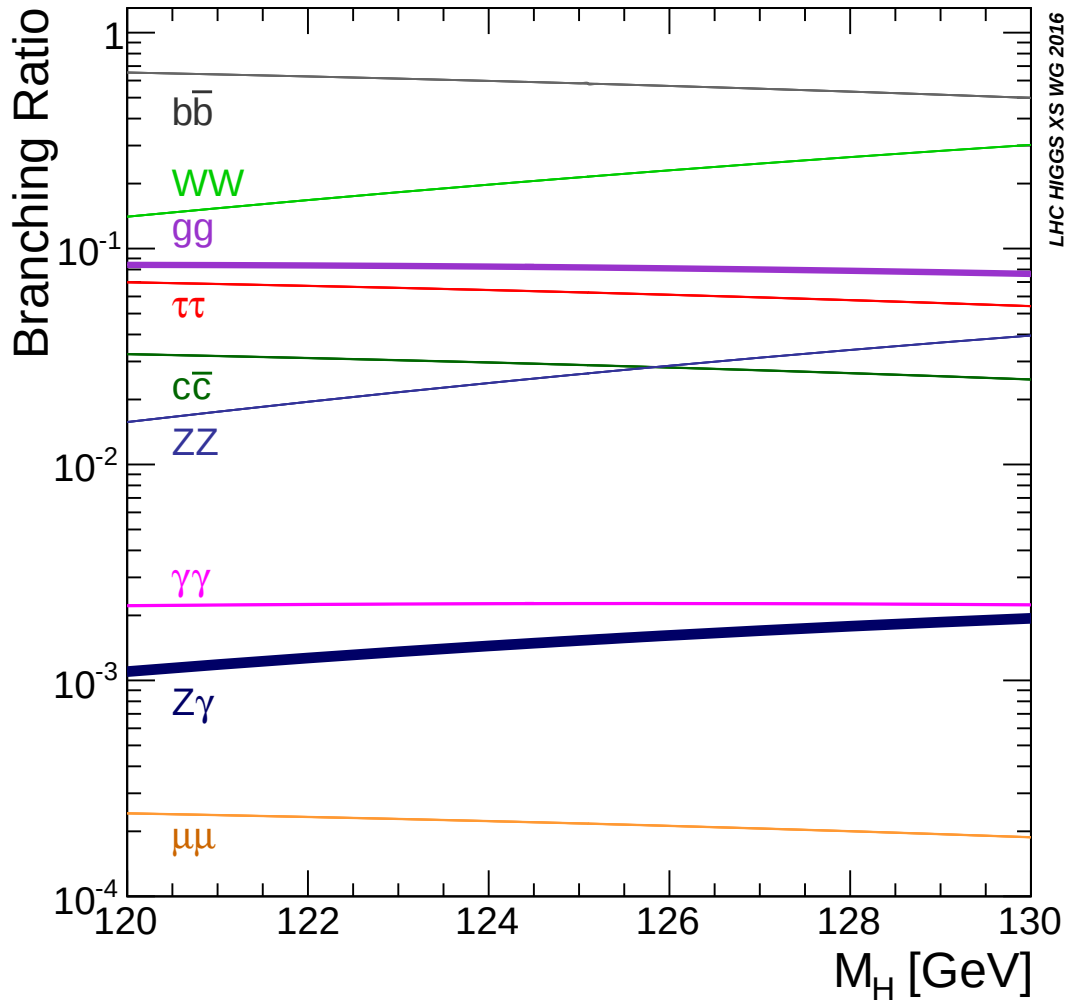


Figure 3.9.: Expected branching fractions of the SM Higgs-boson as function of its mass [26,27]. Branching fraction for $H \rightarrow \tau\tau$ (red line) and $H \rightarrow \mu\mu$ (orange line) are in the order of $\sim 6\%$ and $\sim 0.02\%$, respectively.

3.3. Lepton-flavour violating decays of the Higgs boson at the LHC

In Section 3.1, the natural branching fraction $B_{nat}(H \rightarrow ij)$ was introduced. The LFV branching fractions are expected to be of same size or smaller than $B_{nat}(H \rightarrow ij)$ using the natural assumption. Furthermore, LFV branching fractions $B_{nat}(H \rightarrow ij)$ can be expressed in terms of SM ones. Then, the LFV Higgs boson decay to $\mu\tau$ should have a natural branching fraction in the order of 0.5%, which is between those of decays to taus or muons. Lepton-flavour violating decays of the Higgs boson to $e\tau$ would have a smaller natural branching fraction in the order of $0.001\% = 10^{-5}$. Decays of the Higgs boson to $e\mu$ would even have a smaller natural branching fractions of the order 10^{-6} .

Upper limits on LFV branching fractions Low-energy measurements, described in Section 3.2, indirectly constrain branching fractions of the LFV Higgs-boson decays to $\mu\tau$ or $e\tau$ to be smaller than $\sim 10\%$, which is larger than their natural branching fractions. The constraint for decays to $e\mu$ is in the order of 10^{-8} , which is smaller than the natural branching fraction. These constraints were derived assuming flavour changing neutral currents to be dominated by Higgs boson contributions, therefore lepton-flavour violating effects could be cancelled by effects from other new physics leading to weaker limits [6]. Direct searches for LFV decays of the Higgs boson are independent from assumptions on other new physics models.

A search for lepton-flavour violating decays of the Higgs boson to $\mu\tau$ has been done using the CMS data sample collected in proton-proton collisions at a centre-of-mass energy of $\sqrt{s} = 8$ TeV [17], which was the first direct search for LFV decays of the observed Higgs boson. Constraints on the branching fraction $B(H \rightarrow \mu\tau) < 1.51\%$ could be set. Figure 3.10 shows the expected and observed 95% CL upper limits for each category of the search and their combination. The expected limit for the combination is very close to the natural expectation, assuming similar LFV couplings of the Higgs boson as for the SM Higgs-boson decays. This search has improved the indirect limit on the branching fraction $B(H \rightarrow \mu\tau)$ by one order of magnitude. The search was performed in the $\mu\tau_e$ channel, which has the same signature as the $e\tau_\mu$ channel of a corresponding search for decays of the Higgs boson to $e\tau$, but different kinematics of the leptons. Thus, it is expected that a search for LFV decays of the Higgs boson to $e\tau$ in the $e\tau_\mu$ channel is sensitive for branching fractions up to $\mathcal{O}(1\%)$, which would also improve upon indirect limits by one order of magnitude.

The status presented in this section was the status at the time, when the search for LFV decays of the Higgs boson to $e\tau_\mu$, described in this thesis, was started. Other results from searches of LFV decays of the Higgs boson have been published by the ATLAS and the CMS collaborations. They are discussed in Chapter 9.

3. Lepton-flavour violating decays of the Higgs boson

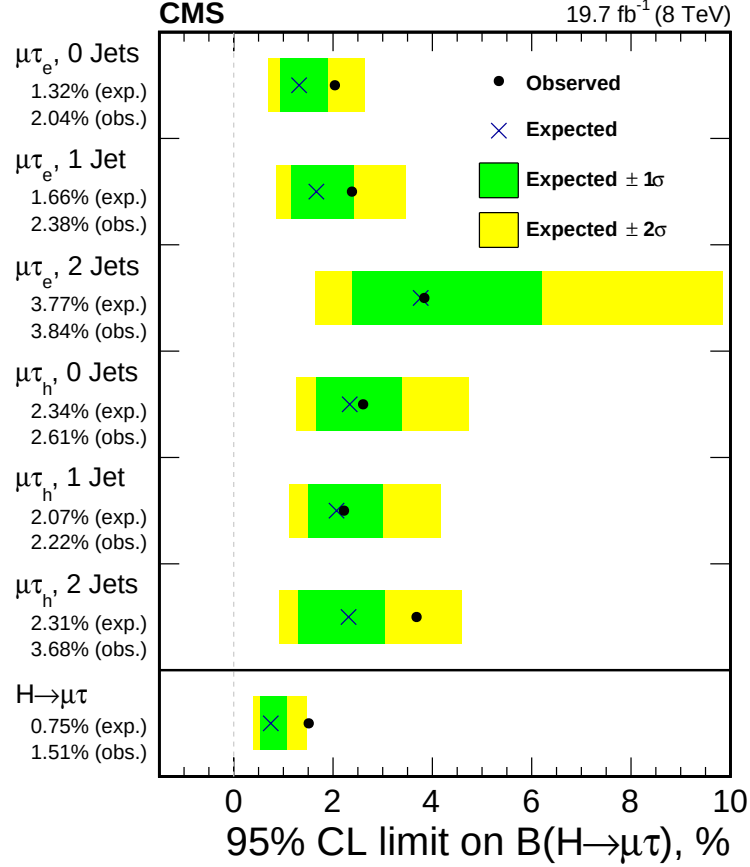


Figure 3.10.: Expected and observed 95% CL upper limits by category of a search for lepton-flavour violating $H \rightarrow \mu\tau$ decays [17]. The data sample used was collected in proton-proton collisions at a centre-of-mass energy of $\sqrt{s} = 8$ TeV with the CMS experiment and corresponds to an integrated luminosity of 19.7 fb⁻¹.

4. Experimental Setup

Exploring the physics at the TeV scale and to study the mechanism of electroweak symmetry breaking including the search for the SM Higgs boson, are prime goals of the Large Hadron Collider (LHC) and the Compact Muon Solenoid (CMS) experiment. Searches for new particles predicted by beyond the SM physics are also main goals of the LHC and CMS physics program. Additionally, studies of the SM top quark properties, electroweak physics, or heavy flavour physics, physics of hadrons containing a charm or bottom quark, are of main interest for the physics program. High energy heavy-ion collisions address the physics of strongly interacting matter and the quark-gluon plasma, a presumed state which exist at extreme values of energy density and temperature and consists of asymptotically free quarks and gluons. Hadrons are consist of quark and gluons, therefore the initial energy of two colliding particles is not known. This is different at lepton colliders, where each particle has the same energy and the energy of the collision is known. Thus, lepton colliders are well suited for precision measurements, while hadron colliders can explore a wide range of collision energies.

First, the Large Hadron Collider [7] and its experiments are introduced with the focus on the Compact Muon Solenoid (CMS) experiment [28]. Their physics programs are discussed and an overview of the data-taking periods is given for the proton-proton collisions. Afterwards, the CMS detector is introduced. This section is based on the technical design report [29]. First, its coordinate system and basic kinematic properties are discussed. Then, an overview of the detector is given followed by an short discussion of each subdetector. The online event selection (trigger) and data acquisition system are also discussed. Finally, luminosity and its measurement are presented [29–33].

4.1. The Large Hadron Collider

The Large Hadron Collider (LHC) [7] is a two-ring-superconducting-hadron accelerator located at CERN near Geneva. It was designed to collide proton beams with a beam energy of 7 TeV leading to a centre-of-mass energy $\sqrt{s}$ of 14 TeV. The design luminosity for proton-proton collisions was $10^{34} \text{ cm}^{-2}\text{s}^{-1}$. Heavy lead (Pb) ions can be accelerated up to an energy of 2.8 TeV per nucleon and a peak luminosity of $10^{27} \text{ cm}^{-2}\text{s}^{-1}$ can be

4. Experimental Setup

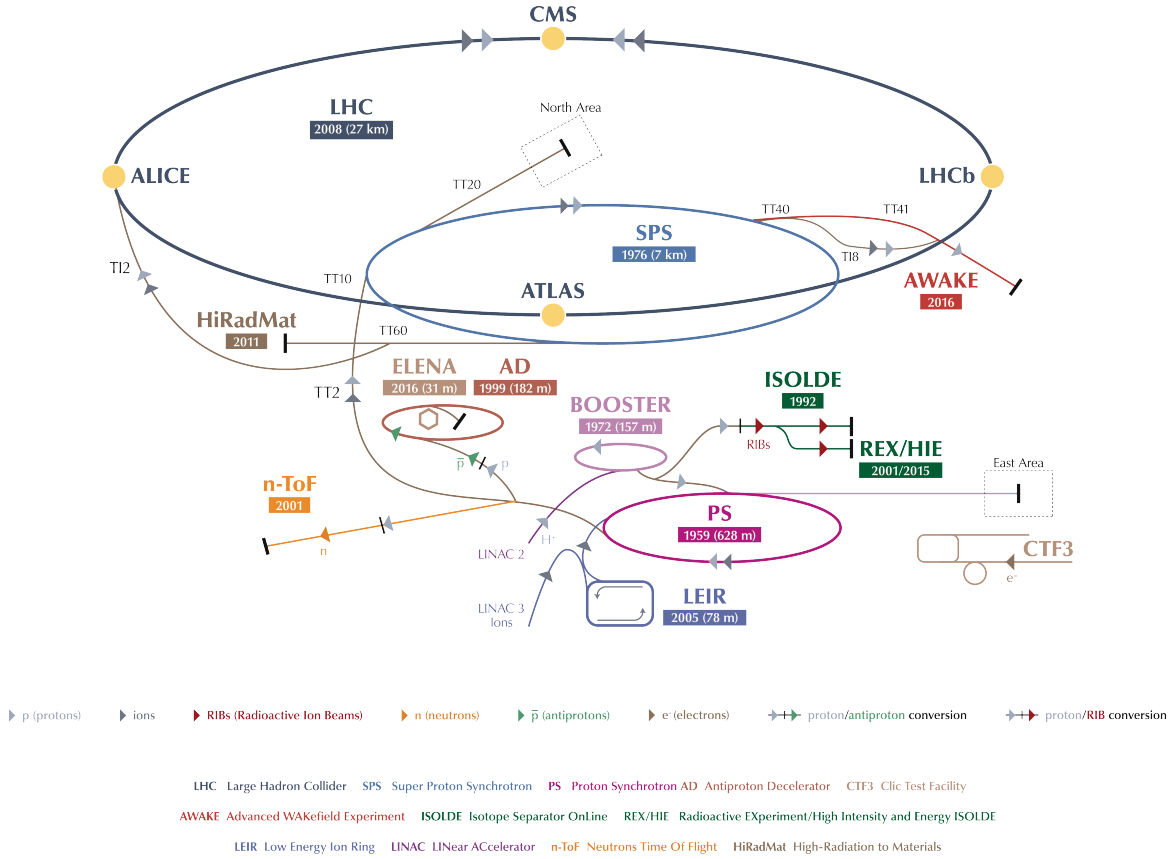


Figure 4.1.: The CERN accelerator complex. Particles are accelerated in a complex chain of smaller particle accelerators to their final energies before they are injected into the LHC (dark blue line) [34].

reached. The LHC with its unprecedented collision energy is the most powerful tool for particle physics research today. Its energy is only outperformed by the energy of cosmic rays. The existing 26.7 km tunnel constructed between 1984 and 1989 for the Large Electron-Positron collider (LEP) [35] was reused for the installation of the LHC. There are eight straight sections and eight arcs lying 45 – 170 m below the surface with a 1.4% slope towards the Lake Geneva. Counter-rotating beams¹ and two rings are used as the LHC is a particle-particle collider, unlike particle-antiparticle colliders that can use a single ring for both beams. As the LHC collides heavier particles than the LEP collider, it suffers less from synchrotron radiation.

The CERN accelerator complex, shown in Figure 4.1, acts as an injector of the protons and heavy ions. Hydrogen gas from a bottle is used as source of the protons. The electrons are stripped off and only the protons enter the Linear accelerator 2 (Lineac2) [36], where they are accelerated to 50 MeV [37]. Then, the protons are further accelerated in the Proton Synchrotron Booster (PSB) [38] to 1.4 GeV. Afterwards, they are injected in the

¹beams with corresponding protons moving in opposite directions



Figure 4.2.: Aerial View of the CERN taken in 2008 [39].

Proton Synchrotron (PS) [40], where they are accelerated to 25 GeV. A bunch train² with the given LHC spacing is produced within the PS before extraction. Finally, the protons are accelerated in the Super Proton Synchrotron (SPS) [41] to 450 GeV before they are injected in the LHC. Further information on the injector chain can be found in [7, 29, 37].

Four interaction points are equipped with particle detectors: A Toroidal LHC Apparatus (ATLAS) experiment [42], the Compact Muon Solenoid (CMS) experiment [28, 29], A Large Ion Collider Experiment (ALICE) [43], and a Large Hadron Collider beauty (LHCb) experiment [44]. Two further smaller experiments, a Total, Elastic and diffractive cross-section Measurement (TOTEM) [45] and the Large Hadron Collider forward (LHCf) [46] experiment are located near the CMS interaction point and near the ATLAS interaction point, respectively. Figure 4.2 shows location of the four larger experiments.

²Groups of proton bunches, where the space between bunches in the train is smaller than between the trains [47].

4. Experimental Setup

The ATLAS and CMS experiments The ATLAS [42] and CMS [28, 29] experiments have both multi-purpose detectors installed at the interaction points. Both detectors were build to detect particles from proton-proton (p-p) or heavy-ion (Pb-Pb) collisions. The search for the Higgs boson was one of their main tasks until a Higgs boson was discovered by both collaborations in 2012 [3, 4, 13]. Henceforth, the main tasks is to study the production and the decay of the discovered Higgs boson, to disentangle its properties and to check if it is the SM Higgs boson or a Higgs boson of an extension of the SM. High precision tests of QCD, electroweak interactions, and heavy flavour physics are further tasks of both experiments. Furthermore, precision measurements of production, the couplings and spin of the top quark are also pursuit. Several searches for physics beyond the SM are also part of their physics program, among them searches for supersymmetric particles and further Higgs bosons.

The ALICE experiment ALICE [43] is a dedicated heavy-ion detector. It focuses on QCD describing the strong interaction. It addresses the quark-gluon plasma, which is a presumed state with asymptotically free quarks and gluons. Particles produced in lead-lead (Pb-Pb) collisions, p-Pb, or p-p collisions are detected by the ALICE detector. The extreme particle multiplicity in nucleus-nucleus collisions, requires a high-momentum resolution as well as an excellent particle identification of hadrons, electrons, muons, and photons. Data taken with proton beams is used as reference data for heavy-ion beams.

The LHCb experiment The LHCb [44] experiment is specialised on heavy flavour physics. It is looking for indirect evidence of new physics in CP violation and rare decays of bottom and charm hadrons. Current results in heavy flavour physics are fully consistent with the SM. On the other hand, the SM weak interaction cannot explain the large asymmetry between the amount of matter and anti-matter in the universe. Thus, a new source of CP violation beyond the Standard Model is needed. Many models of new physics change the expectations of the CP-violating phase, the rare decay branching fractions and allow modes which are forbidden in the SM. The experiment runs at lower luminosities in order to reduce the number of interactions per bunch crossing by changing the beam focus at its interaction point independently from the other interaction points. This leads to a lower occupancy and also reduces the radiation damage. An excellent vertex and momentum resolution is needed for reconstruction of the bottom-hadron decays. Furthermore, a good proper-time resolution is needed for studying the oscillating $B_s - \bar{B}_s$ meson system, where the B_s consists of a strange and an anti-bottom quark [1]. Additionally, an excellent identification of electrons, muons, photons, protons, pions, η mesons and kaons [1] is needed for studying the different bottom hadron decays modes.

The TOTEM experiment TOTEM [45] measures the total proton-proton cross section with a luminosity-independent method. Elastic and diffractive scattering at small angles are studied for a deeper understanding of the proton structure. TOTEM is an independent detector but it is technically integrated into CMS. A good acceptance for particles produced at very small angles with respect to the beam is required. Two tracking telescopes are installed on each side in the pseudorapidity region $3.1 \leq |\eta| \leq 6.5$, where pseudorapidity η is defined in Section 4.2.1. Additionally, very forward charged particles are detected in movable beam-pipe insertions, which are called Roman Pots.

The LHCf experiment The LHCf [46] experiment measures neutral particles emitted in the very forward region of the LHC collisions. Data is provided for calibrating the hadron interaction models that are important to study extremely high-energy cosmic-rays. Highest energy cosmic-rays with energies above 10^{19} eV have a great scientific interest because their origin, propagation and interactions are unknown. Understanding the results of cosmic-ray experiments strongly depends on simulations. Perturbative QCD models used for simulation of air shower experiments are phenomenological models which have not been verified for energies close to the 10^{20} eV. Thus, experimental data is needed to calibrate these models. The design centre-of-mass energy of the LHC, $\sqrt{s} = 14$ TeV, corresponds to a laboratory equivalent energy up to 10^{17} eV. The total inelastic cross section and the particle production energy spectra at very forward angles determine the development of air showers. LHCf focuses on the second quantity, while TOTEM is specialised for the first quantity. Two independent shower calorimeters made of a pair of sampling and imaging calorimeters identify γ -rays, measure their energy spectrum, measure the incident position and reconstruct the γ - γ invariant mass distribution for reconstruction of neutral pions.

LHC data-taking periods Not only is the LHC operated using proton-proton collisions or heavy-ion collisions, but it is also operated at different centre-of-mass energy using different spacing between the proton bunches. The LHC was operated at $\sqrt{s} = 7$ TeV in 2010 and 2011 during LHC Run I data-taking. Additionally, it was operated at $\sqrt{s} = 8$ TeV in 2012. A *bunch spacing* of 50 ns was used in LHC Run I. In 2013, there was a long shutdown for preparing the LHC to run at the higher energy of $\sqrt{s} = 13$ TeV. In the LHC Run II data-taking, a bunch spacing of 25 ns was used and proton-proton collision data were collected at $\sqrt{s} = 13$ TeV in 2015 and 2016. Data from heavy-ion collisions were also collected, but this data is not used in this thesis. The integrated luminosity delivered to CMS as function of time is shown in Figure 4.3 for each proton-proton collision data-taking period. Not the whole delivered data is recorded by CMS and only the part of it considered as good is used for physics analysis. Section 4.3 discusses how the luminosity is measured. Data used in the search for lepton-flavour violating decays of the Higgs boson described

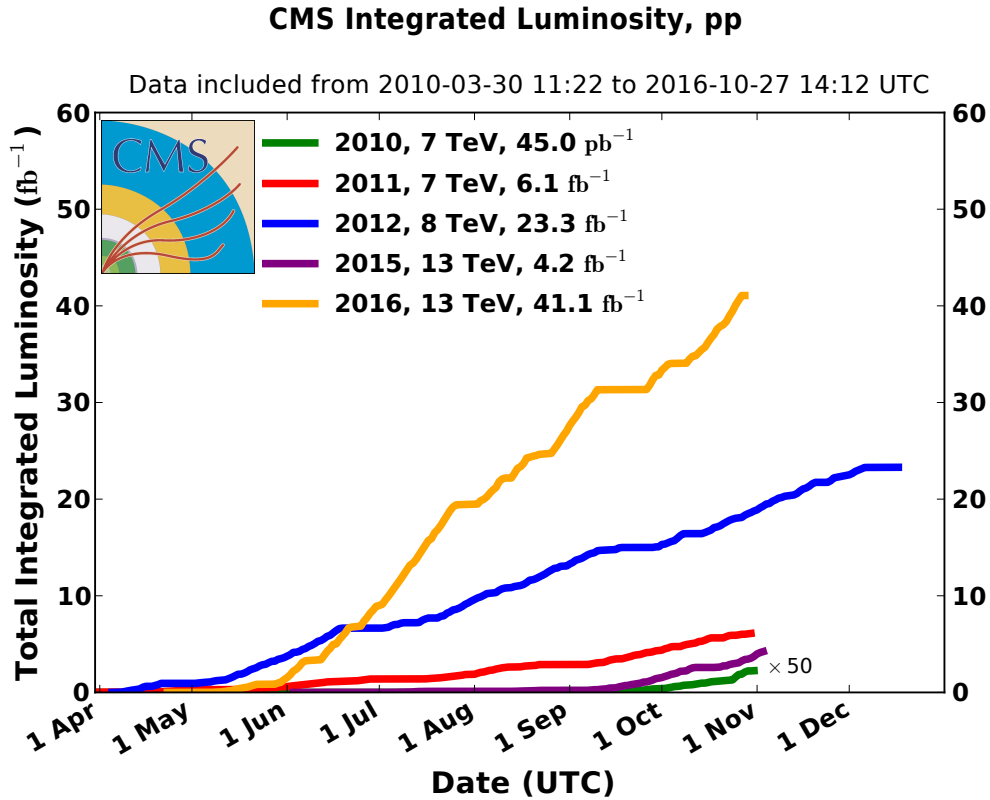


Figure 4.3.: Cumulative luminosity as function of time delivered to CMS during stable beams for proton-proton collisions at different centre-of-mass energies. The luminosity is shown for LHC Run I data-taking in 2010 (green), 2011 (red), and 2012 (blue). It is also shown for LHC Run II data-taking in 2015 (purple) and 2016 (orange) [48].

in this thesis was collected in proton-proton collisions at $\sqrt{s} = 8$ TeV and $\sqrt{s} = 13$ TeV. Table 4.1 gives a summary of the data used in this search.

Table 4.1.: Certified data collected in proton-proton collisions used in the search for lepton-flavour violating decays of the Higgs boson.

Year	$\sqrt{s}$ [TeV]	Run period	Integrated luminosity
2012	8	LHC Run I	19.7 fb ⁻¹
2015	13	LHC Run II	2.3 fb ⁻¹

4.2. The Compact Muon Solenoid experiment

The Compact Muon Solenoid (CMS) is a multipurpose detector [28, 29]. One of the prime goals of CMS is to study the mechanism of electroweak symmetry breaking.

A huge step was achieved by the discovery of the Higgs boson in 2012 by the ATLAS and CMS experiments [3, 4]. After the discovery, precision measurements of the properties of the Higgs boson are one of the main tasks of the experiment. Another prime goal of CMS is to study the physics at the TeV scale. On the one hand, to study the SM processes, such as QCD, b-quark physics, top-quark properties, and electroweak physics like W and Z boson properties, on the other hand searches for new physics are also important, among them searches for supersymmetric partners of the SM particles, searches for extra dimensions, new vector boson high

mass states and many other new physics models. The search for the Higgs boson was done over a wide range of decay modes depending on the mass of the Higgs boson. Final states that contain isolated leptons and photons have a clean signature, therefore searches for the Higgs boson were done using signatures involving leptons and photons. *Supersymmetry* (SUSY) is a generalisation of space-time symmetries of QFT which transforms fermions into bosons and vice versa [1]. SM particles have supersymmetric partners which differ in spin by half a unit. Most SUSY models investigated by CMS require that *R-parity* is conserved, which relates the baryon and leptons numbers and spin. SM particles have even *R-parity*, while SUSY particles have odd *R-parity*. Consequently, SUSY particles are produced in pairs if *R-parity* is conserved. Furthermore, the lightest supersymmetric particle (LSP) is stable and does not interact with the CMS detector leading to an imbalance of the observed total momentum. Decay of SUSY particles lead to signatures with leptons, bunches of hadrons originating from hadronisation of quarks and gluons, and hard isolated photons in the detector. Searches for new massive vector bosons require a good momentum resolution at high momenta. Good identification of hadrons originating from b-quarks are needed for top-quark physics and heavy-flavour physics. Heavy ion physics requires a large detector acceptance, high-resolution calorimeters and tracking devices, and flexible triggers.

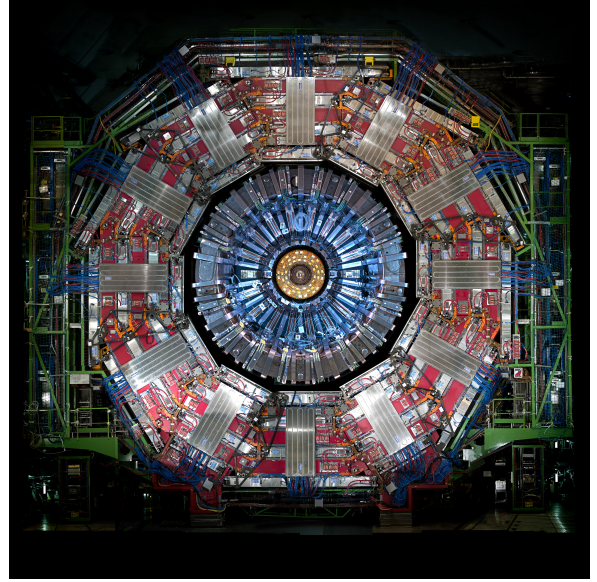


Figure 4.4.: The Compact Muon Solenoid experiment [49].

4. Experimental Setup

4.2.1. The coordinate system of CMS

CMS uses a right-handed Cartesian coordinate system with its origin at the centre of the detector, which is the nominal interaction point. The x -axis points towards the centre of the LHC, while the y -axis points upward towards the surface and perpendicular to the plane of the LHC ring. Thus, the z -axis points along the anticlockwise beam-direction. Two angles are defined, where the azimuthal angle ϕ is measured from the x -axis in the x - y plane and the polar angle θ is measured from the z -axis. Pseudorapidity is defined as

$$\eta = -\ln \left[\tan \left(\frac{\theta}{2} \right) \right] \quad (4.1)$$

and describes the angle of a particle relative to the beam axis, illustrated in Figure 4.5. Distances in ϕ and η are denoted $\Delta\phi$ and $\Delta\eta$. These distance measures are used to define cones around an axis with a border defined by ΔR with

$$\Delta R = \sqrt{(\Delta\phi)^2 + (\Delta\eta)^2}. \quad (4.2)$$

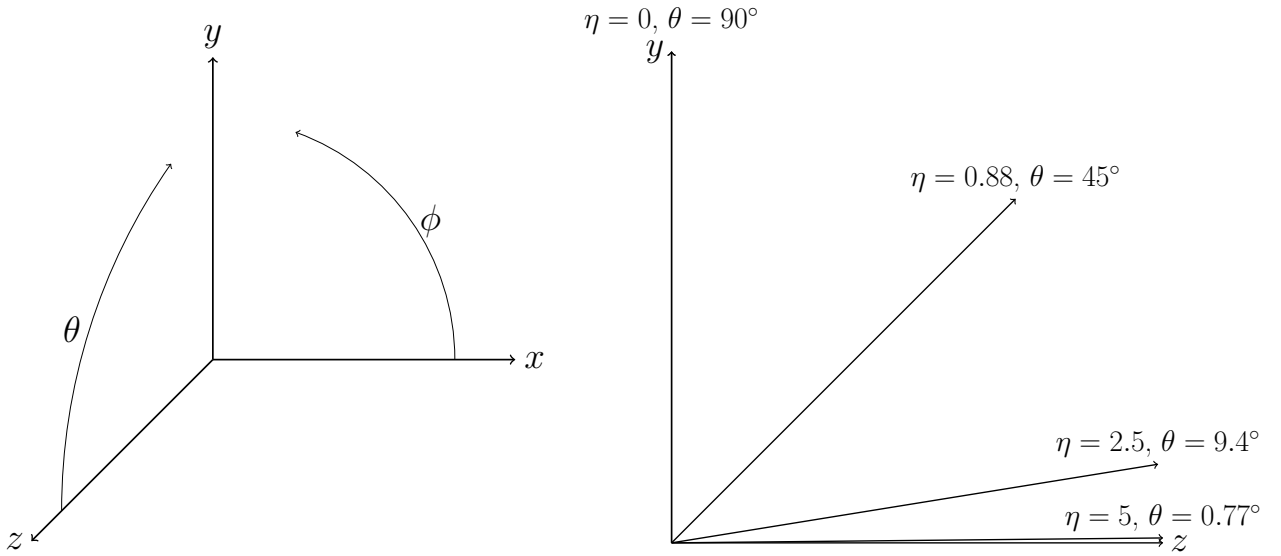


Figure 4.5.: Coordinate system convention of CMS (left) and the relation between pseudorapidity η and polar angle θ (right).

4.2.2. Kinematic quantities

At proton-proton colliders, such as the LHC, the protons carry half of the collision energy $\sqrt{s}$. The hard interactions does not take place between the colliding protons but between two of their partons. Each parton carries a fraction x_i of the proton momentum. As the parton mass can be neglected with respect to the momentum $\vec{p}$, the parton energy is given

by its momentum. In the laboratory frame, the four-momenta of the partons are

$$\begin{aligned} p_1 &= x_1 \cdot \frac{\sqrt{s}}{2}(1, 0, 0, 1), \\ p_2 &= x_2 \cdot \frac{\sqrt{s}}{2}(1, 0, 0, -1). \end{aligned}$$

Then, the invariant mass M of the hard collision is given by

$$\hat{s} \equiv M^2 = (p_1 + p_2)^2 = \frac{s}{4} \cdot [(x_1 + x_2)^2 - (x_1 - x_2)^2] = x_1 x_2 s,$$

where $\sqrt{\hat{s}}$ denotes the centre-of-mass energy of the parton-parton collision. The hard collision products have a total momentum which is zero for the x - and y -components but in general non-zero for the z -component. Thus, the momentum and energy of the hard-collision products is measured transverse to the beam direction in the x - y plane. The transverse momentum and the transverse energy are denoted p_T and E_T , respectively, with $E_T = E \cdot \sin \theta$.

Only weakly interacting particles, such as neutrinos, do not produce a signal in CMS detector. However, such particles have a certain signature. They lead to an imbalance in the observed total transverse momentum in the event, *missing transverse energy*, denoted as $\cancel{E}_T$.

4.2.3. Detector requirements and layout

Final states that contain isolated leptons and photons leave clean signatures in the CMS detector. A good muon identification and momentum resolution is needed over a wide range of momenta. The muon charge has to be determined unambiguously. Furthermore, a good momentum resolution and reconstruction efficiency of charged particles is crucial for the reconstruction of electrons and charged hadrons. Charged-particle momenta are measured using the curvature of their trajectory. A large bending power is needed to measure precisely the momentum of particles with large momentum. Additionally, a good electromagnetic energy resolution is needed and a wide geometrical coverage. The direction of the photons and/or correct localisation of the primary interaction should be measurable. A large background of photon identification arise from neutral mesons, which are collimated and are reconstructed as a single photon [50]. This background is dominated by π^0 mesons which decay almost solely into two photons. These π^0 mesons have to be distinguished from isolated photons and leptons, also at high luminosities.

The search for the Higgs boson decay into two taus is important as the coupling of the SM Higgs boson to taus is larger than for the other leptons. Furthermore, the Higgs decay to two b -quarks has the largest branching fraction for Higgs masses up to ~ 135 GeV [4].

4. Experimental Setup

Top-quarks decay predominantly to a b-quark and a W boson. Quarks hadronise and the hadrons can be detected with the CMS detector. A good identification of hadronically decaying taus and hadrons from b-quark decays is needed for the Higgs boson and top-quark precision physics. The triggering and identification efficiency of hadronically decaying taus and hadrons from b-quark decays can be improved by a good measurement of the impact parameter of charged-particle tracks and by a good position measurement of the secondary vertices. This requires pixel detectors close to the interaction region.

Hadronic calorimeters with a large hermetic coverage ($|\eta| < 5$) and with a fine lateral segmentation ($\Delta\eta \times \Delta\phi < 0.1 \times 0.1$) are required for a good energy measurement of hadrons and a good estimation of $\cancel{E}_T$.

There are also experimental challenges. At the design centre-of-mass energy and design luminosity, approximately one billion inelastic events per second are expected. Not all of these events can be stored, therefore only interesting events have to be selected. The online event selection process, *trigger*, must reduce the one billion interactions per second to no more than a few hundred events per seconds. Also the short time between two bunch crossings, 25 ns, has a major implication on the design of the readout and trigger system.

Multiple proton-proton interactions (*pileup*) happen in one bunch crossing. The products of the interactions overlap and can be wrongly linked. Long response times of detector elements and their electronic signal longer than 25 ns increase this effect. High granularity detectors with good time resolution result in a low occupancy. This requires a large number of detector channels and therefore, a good synchronisation of the detector electronic channels. The large flux of the particles and resulting high radiation levels require radiation-hard detectors and front-end electronics.

The layout of the CMS detector is shown in Figure 4.6. As the name "Compact Muon Solenoid" indicates, a superconducting solenoid is the heart of CMS. Large bending power for the momentum measurement of charged particles within a compact design is achieved by using a high magnetic field of 3.8 T. The internal part of the magnetic coil is large enough to accommodate the inner tracking system and the calorimetry. Charged-particle trajectories are measured up to $|\eta| < 2.5$ in the inner tracking system, which is composed of a pixel detector close to the interaction region and a silicon strip tracker. The pixel detector can resolve individual vertices and distinguish between vertices from the primary interaction and secondary vertices from the decay of the particles produced in the primary interaction. Trajectories are precisely measured with the high granular silicon strip tracker, which can deal with high charge particle multiplicities.

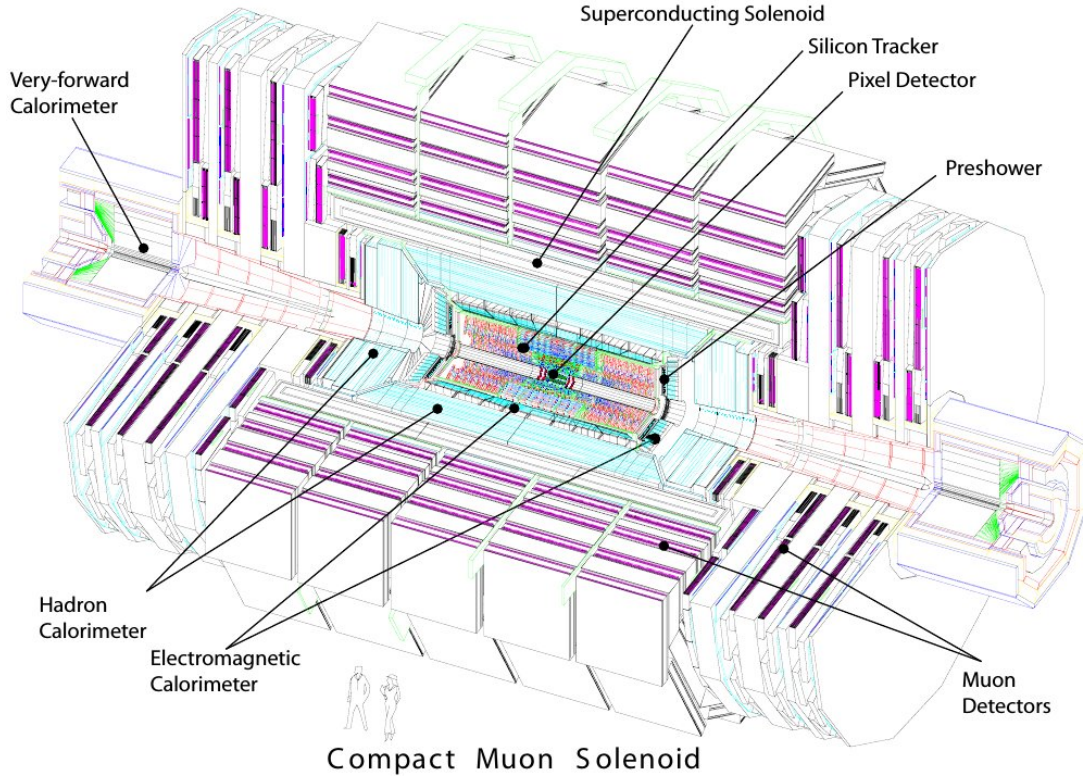


Figure 4.6.: The CMS detector [29].

The energy of the particles is measured in calorimeters [1]. Calorimeter material initiates electromagnetic (EM) or hadronic showers. For electromagnetic interactions, the characteristic interaction length is the radiation length X_0 , while the characteristic interaction length for hadronic showers is the nuclear interaction length λ_I . In *homogeneous calorimeters*, the entire volume is sensitive, while *sampling calorimeters* consist of metallic absorber sandwiched or threaded with an active material which generates the signal. CMS has an EM calorimeter (ECAL) made of lead tungstate PbWO_4 in front of a brass/scintillator sampling hadron calorimeter (HCAL). An additional layer of scintillators is outside the coil. The magnet is used as absorber material. This iron/quartz-fibre calorimeter is referred to as the hadron outer (HO) detector. The return field of the solenoid is large enough to saturate 1.5 m of iron. This allows the installation of up to four muon stations outside the solenoid. The muon detectors are sandwiched between the layers of the steel return yoke. Their main task is the triggering on muons, identifying muons and improving the momentum measurement of muons with high momenta.

The main features of the CMS detector are the high-field solenoid and the full silicon-based inner tracking system. CMS has a total length of 21.6 m and a diameter of 14.6 m. The thickness of the ECAL in radiation lengths is greater than $25 X_0$ and for the HCAL it is greater than $7 - 11 \lambda_I$ depending on η . CMS has a barrel-shaped part in the central region at smaller η termed *barrel* and two forward regions at higher η termed *endcap*.

4. Experimental Setup

4.2.4. Magnet

Charged particle momenta and their charge sign are determined by the curvature of the particle trajectory in a magnetic field. The choice of magnetic field configuration is the most important aspect on muon measurement. A good momentum resolution of $\Delta p/p \approx 10\%$ is required for muons with a momentum of 1 TeV. CMS uses a large superconducting solenoid with a magnetic field of 3.8 T. This is about 100,000 times stronger than the magnetic field of the earth [51]. The solenoid is 13 m long and has an inner diameter of 5.9 m. A high-purity aluminium-stabilised conductor with a four-layer winding is used, which has to withstand an outward pressure. Niobium titanium is used as superconductor material [52]. The conductor is composed out of five coils and its superconducting wire is cooled using an indirect cooling by thermosyphon.

4.2.5. Inner Tracking system

Trajectories of particles are measured in the inner tracking system, which covers the region $|\eta| < 2.5$. Close to the interaction region, the particle flux is the highest. There the silicon pixel detector is placed, allowing the reconstruction of vertices from heavy-flavour hadrons with a b or c quark. Each pixel has a size of $\approx 100 \times 150 \mu\text{m}^2$. The spatial resolution of the radius and azimuthal angle r - ϕ measurement is about $10 \mu\text{m}$, and $20 \mu\text{m}$ for the z -coordinate measurement. In the intermediate region ($20 \text{ cm} < r < 55 \text{ cm}$) and outermost region ($r > 55 \text{ cm}$) silicon microstrips are used with a size of $10 \text{ cm} \times 80 \mu\text{m}$ (minimum cell size) and a size of $25 \text{ cm} \times 180 \mu\text{m}$ (maximum cell size), which provide the required granularity.

The tracking volume in the barrel has the shape of a cylinder with a length of 5.8 m and a diameter of 2.6 m. Figure 4.7 illustrates the inner tracking system of CMS. In the barrel, the pixel detector consist of three layers of at radii of 4, 7, and 11 cm. Additionally, there are 10 layers of silicon microstrip detectors. The strip tracker is divided into a *Tracker Inner Barrel* (TIB) made out of four layers and a *Tracker Outer Barrel* (TOB) made out of six layers. Each part has two layers which provide a single-point measurement in the r - ϕ and r - z coordinates. In the TIB, the single-point resolution is $23 - 34 \mu\text{m}$ in the r - ϕ direction and $23 \mu\text{m}$ in z , while it is $35 - 52 \mu\text{m}$ in the r - ϕ direction and $52 \mu\text{m}$ in z for the TOB. In the two endcaps, there are just two pixel layers and nine microstrip layers. The endcaps of the strip tracker are divided into the *Tracker End Cap* made of 9 disks and the *Tracker Inner Disks* (TID), which are made of three small disks, which fill the gap between the TIB and the TEC.

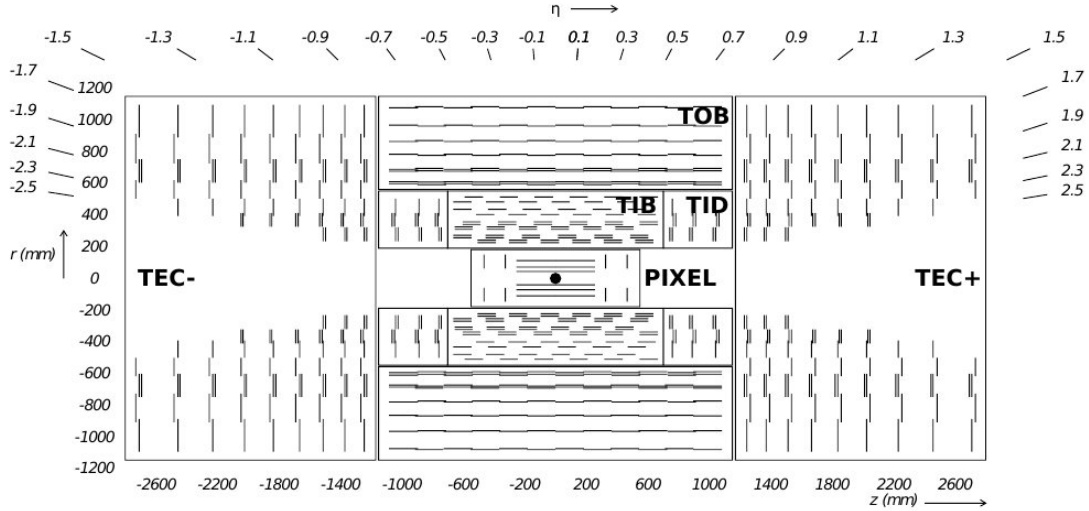


Figure 4.7.: Inner tracking system of CMS [53].

4.2.6. Electromagnetic Calorimeter

The electromagnetic calorimeter (ECAL) measures the energy of electrons and photons. A hermetic, homogeneous crystal (ECAL) with a coverage of $|\eta| < 3$ is used, which has an excellent energy resolution and high granularity, therefore it also has an excellent separation of close clusters [54, 55]. It is used to identify photons, which leave no signal in the inner tracking detector. Lead tungstate (PbWO_4) crystals are used, which have a short radiation length of $X_0 = 0.89$ cm, are radiation hard, and emit blue-green scintillation light with a maximum at 420 nm. The scintillation decay time is short and in the same order as the LHC bunch crossing time. About 80% of the light is emitted within 25 ns. However, the light output depends on the temperature, therefore a cooling system is needed to preserve the energy resolution. Furthermore, a low light yield requires the use of photodetectors with intrinsic gain that can additionally operate in a high magnetic field. The crystals allow a compact calorimeter inside the solenoid that is fast, has a fine granularity and is radiation resistant. Figure 4.8 illustrates the layout of the CMS ECAL. The calorimeter is divided into a barrel ECAL (EB) which covers the region $0 < |\eta| < 1.479$ and two ECAL endcaps (EE) which cover the region $1.479 < |\eta| < 3.0$. Preshower detectors (ES) are installed in front of each endcap and cover the region $1.653 < |\eta| < 2.6$.

In the EB, silicon avalanche photodiodes (APDs) are installed to detect the scintillation light. They also respond to temperature changes, therefore they require a stable temperature. The barrel section is structured in 36 identical supermodules, 18 in each half barrel, with the crystals being arranged in a η - ϕ grid and covering $\Delta\eta \times \Delta\phi = 0.0174 \times 0.0174$. Vacuum phototriodes (VPTs) are installed in both EE for detecting the scintillation light. The crystals are arranged in units of 5×5 crystals, termed *supercrystals* (SC). Both, the crystals and the SCs, are arranged in a rectangular x - y grid. Neutral pions dominantly

4. Experimental Setup

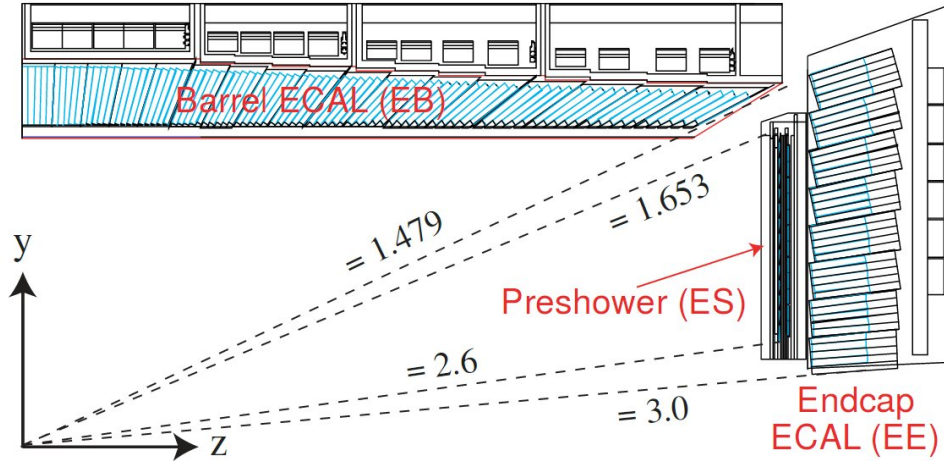


Figure 4.8.: Layout of the CMS electromagnetic calorimeter [29].

decay to two photons, where two closely separated photons can mimic high-energy photons, therefore a preshower system is installed in front of each EE in order to identify and reject π^0 mesons. It also improves the position determination of electrons and photons as it has a higher granularity than the EE. The preshower detector is a sampling calorimeter consisting of two layers of lead absorber which initiate the electromagnetic showers from incoming electrons or photons. Each lead radiator is followed by silicon strip sensors which measure the energy deposits and the transverse shower profiles. The strips in both planes of silicon sensors are orthogonal oriented and are placed after a radiation length of $2 X_0$ and $1 X_0$ of the lead absorber.

4.2.7. Hadron Calorimeter

Energy of hadrons is measured in the hadron calorimeter (HCAL). Most of the HCAL calorimetry is located inside the magnetic coil surrounding the ECAL system. An important requirement on the HCAL design is to minimise the non-Gaussian tails in the energy resolution and to provide a good hermeticity for the determination of $\cancel{E}_T$. It is designed to maximise the interaction length of the material within the magnet coil. On the other hand, the amount of space devoted to the active medium is minimised.

The HCAL is a sampling calorimeter which is divided into a hadron barrel (HB) detector, covering the region $|\eta| < 1.4$ and two hadron endcap (HE) detectors, covering $1.3 < |\eta| < 3.0$. Brass³ is used as absorber material. Brass has a reasonably short interaction length, is easy to machine and it is non-magnetic. A tile/fibre technology made of plastic scintillator tiles is used as active medium. For the innermost and the outermost layer, stainless steel is used for structural strength. The 3.7-mm-thick scintillator plates are sandwiched between the absorber plates. Wavelength-shifting (WLS) fibres, embedded to

³Cartridge brass (C260): 70% copper and 30% zinc

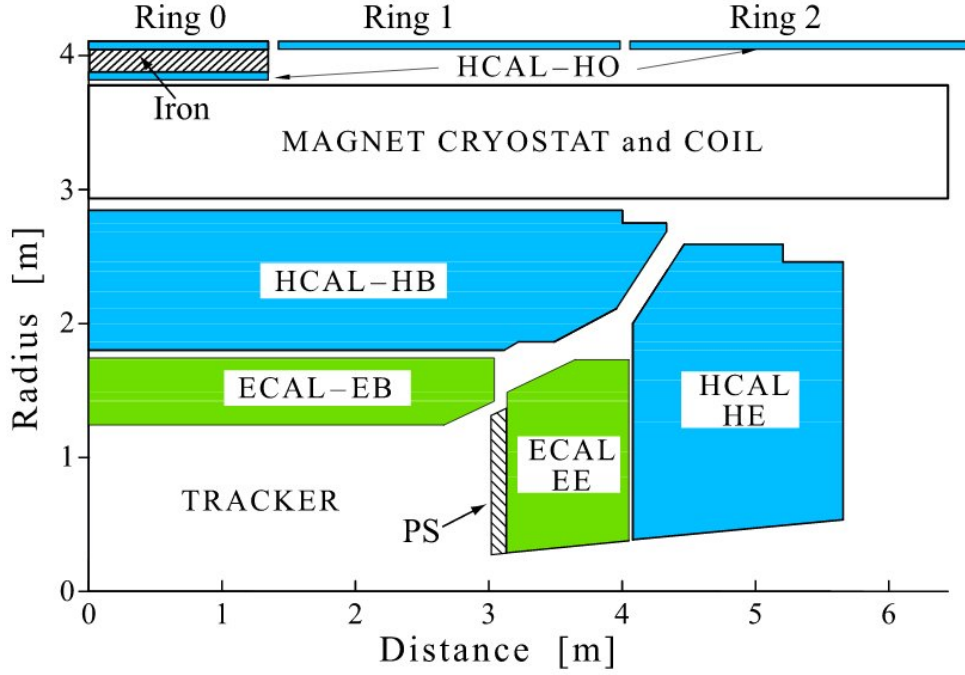


Figure 4.9.: Location of the ECAL and HCAL detectors in and around the CMS magnet [56].

the scintillator tiles, convert the scintillation light. Then, it is channelled to photodetectors via optical fibres. Multi-channel hybrid photodiodes (HPDs) detect the light. They can operate in high axial magnetic fields. The HB is read out as a single longitudinal sampling with a segmentation $\Delta\eta \times \Delta\phi = 0.087 \times 0.087 \approx 5^\circ \times 5^\circ$, which are termed *tower*. In the HE, the ϕ segmentation is $5^\circ - 10^\circ$ and the η segmentation is $0.087 - 0.35$ depending on η . There are essentially no uninstrumented cracks or dead areas in ϕ for the barrel and endcap HCAL.

Additional layers of scintillators are installed outside the coil within the return yoke using the iron as absorber. They sample energy from hadron showers, which leak through the rear of the calorimeters. The central shower containment and the $\cancel{E}_T$ resolution of the calorimeter are thus improved. This sampling calorimeter is referred to as Hadron Outer (HO) detector. It is located along the barrel muon system, therefore its segmentation closely follows that of the barrel muon system. It is divided into five sections along η termed *rings*. The HO follows the HCAL barrel geometry in η and ϕ and covers the region $|\eta| < 1.26$.

Figure 4.9 shows the location of the HB, HE, and HO, respectively. Two further detectors, which cover the region $2.9 < |\eta| < 5.0$, are not shown. They are specialised to measure energetic forward hadronic showers and ensure full geometric coverage for the transverse energy measurement. The Hadron Forward (HF) sampling calorimeters use steel as absorber material and quartz fibres as active medium, which run parallel to the beam. Cherenkov light is emitted by shower particles in the quartz fibres. This light is

4. Experimental Setup

channelled by the fibres to photomultipliers. There are two different types of quartz fibres, long ones (1.65 m) and short ones (1.43 m). Neutral components of the hadron showers are preferentially sampled in the HF, leading to narrower and shorter hadronic showers. This is ideally suited for the forward region. The towers in the HF have a η segmentation of $0.1 - 0.3$ and a ϕ segmentation of 10° , except for the highest η -towers, where the ϕ segmentation is 20° .

4.2.8. Muon System

The Muon system is installed in the magnet return yokes of CMS. Its main tasks are the identification of muons, improving the p_T measurement and charge-sign determination of high- p_T muons. Additionally, it is used to trigger on muons. It is divided into a barrel detector (MB) covering $|\eta| < 1.2$ and two endcap detectors (ME), covering the region $0.9 < |\eta| < 2.4$, defining three regions: the barrel region ($|\eta| < 0.9$), the overlap region ($0.9 < |\eta| < 1.2$), and the endcap region ($1.2 < |\eta| < 2.4$).

Three different types of gaseous detectors are used in the different radiation environments. Drift tube (DT) chambers are used in the barrel region, where the muon rate, the neutron induced background rate, and the residual magnetic field are low. In the two endcaps, the muon rate, neutron induced background, and the magnetic field are high, therefore cathode strip chambers (CSC) are used in this region. Resistive plate chambers (RPC) are used in both subdetectors covering the region $|\eta| < 1.6$. They are operated in avalanche mode to ensure good operation at high rates. The RPCs have double gaps with a width of 2 mm filled with gas. They have a fast response with a good time resolution but a coarser position resolution than the DTs and CSCs, therefore they are used to identify the correct bunch crossing. Figure 4.10 gives an overview of the layout of the muon system.

The MB is divided into 4 stations arranged in cylinders interleaved with the iron yoke. Additionally, it is divided in five wheels along the beam direction following the five wheels of the return yokes. Each chamber consists of 12 layers divided into 3 *Super Layers* (SL), which are made out of four layers of DTs. Two SL measure the r - ϕ coordinate, while a third SL sandwiched in between them measures the z coordinate. In the last muon station, there are only the two SL to measure the r - ϕ coordinate. Each DT chamber has one or two associated RPCs. The single point resolution of the DTs is $\approx 200 \mu\text{m}$. In each endcap, the CSCs and RPCs are arranged in four disks. They are divided into three or two concentric rings in the innermost station and in the other stations, respectively. Each CSC measures up to six space coordinates (r , ϕ , z) and the provided spatial resolution is $\approx 200 \mu\text{m}$, except for the innermost ring in the first disk, where it is $\approx 100 \mu\text{m}$. The angular resolution in ϕ is $\sim 10 \text{ mrad}$. Two independent and complementary sources of information come from the DTs or CSCs and the RPCs, which feed the trigger system.

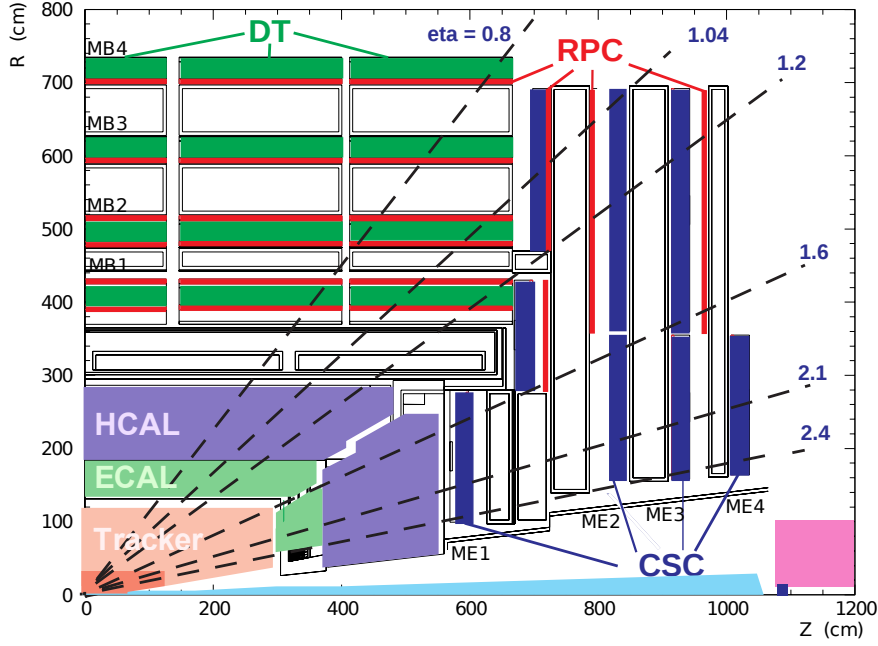


Figure 4.10.: Layout of the CMS muon system [29].

4.2.9. Trigger and data acquisition system

At design luminosity of $10^{34} \text{ cm}^{-2} \text{ s}^{-1}$ and 25 ns time between two bunch crossings, the LHC bunch crossing rate is 40 MHz and about 10^9 interactions take place per second. Out of these interactions, only of the order of 100 interactions per second can be written to archival media. Thus, an online event selection process (trigger) with a rejection factor of 10^6 is needed to reduce the data volume. The main task of the trigger system is to reduce the rate of events while keeping a high efficiency for interesting physics events. Furthermore, the simulation of events must be able to reproduce the sample composition and the correct trigger efficiency.

The trigger and data acquisition system of the CMS experiment is divided into four parts, summarised in Figure 4.11. First, the detector electronics detects signals and pass them to the *Level-1 trigger* processors. The Level-1 trigger logic is located in a service cavern and selects events of interest. Due to the size of the CMS detector and the underground caverns, the transit and time needed for a decision is about $3.2 \mu\text{s}$. However, the time needed for Level-1 trigger calculations is only $1 \mu\text{s}$. In this time, the data is held in pipelined memory buffers. Custom hardware processors form the Level-1 decision to keep or to discard data using the calorimetry, the muon systems, and correlation information between them. *Trigger primitive* objects, such as electrons or muons above a set E_T or p_T thresholds are the basis of the decision taking. Reduced granularity and resolution data are used to form the trigger objects. Sums of E_T and $\cancel{E}_T$ are also employed. Only about one out of 1000 interactions is retained. If the data of an interaction is kept,

4. Experimental Setup

the data is transferred after fixed time interval of about $3.2 \mu\text{s}$ to the front-end readout buffers. The data is further processed and compressed and placed in dual-port memories. A processor farm is responsible for filtering the events. Each processor runs the same *High-Level Trigger* (HLT) code, reducing the output rate from a few kHz to a few hundred Hz for mass storage. Using a processor farm for event filtering allows to benefit from evolution of computer technology and maximises the flexibility in selection of data and of the algorithms. Uninteresting events are discarded as soon as possible. Whenever possible, only necessary objects are reconstructed and only needed information of some subdetectors are used. The idea of partial reconstruction and of many virtual trigger levels is the basis for the HLT. First, calorimeter and muon information is used, followed by the use of pixel tracker data. Finally the full event information is used including full tracking. Identification and isolation criteria as well as minimal energy or transverse momentum thresholds are used by the HLT system [55].

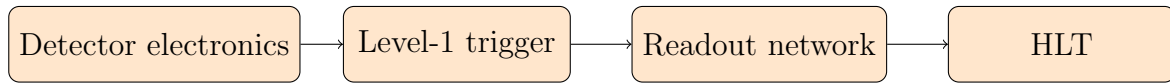


Figure 4.11.: Overview of the trigger and data acquisition system

4.3. Luminosity measurement

Luminosity $\mathcal{L}$ is defined as the ratio of the event rate $\dot{N}$ to the cross section σ of a given process, which is the effective area quantifying the likelihood of a scattering event.

$$\mathcal{L} = \frac{\dot{N}}{\sigma} \quad (4.3)$$

The cross section is measured in units of area, thus the luminosity is given in units of events per time per area, $(\text{b} \cdot \text{s})^{-1} = 10^{24} \text{ cm}^{-2}\text{s}^{-1}$. Then, the *integrated luminosity* L is the integral over the instantaneous luminosity:

$$L = \int \mathcal{L}(t) dt. \quad (4.4)$$

For a given process, the number of expected events which are produced is given by the product of the integrated luminosity and the production cross-section σ_{exp} :

$$N_{\text{exp}} = L \cdot \sigma_{\text{exp}}. \quad (4.5)$$

Thus, the integrated luminosity has to be known in order to estimate the number of events for a given process or to measure the production cross section for a process. Luminosity measurements are used to monitor the LHC performance in real time and they provide an overall normalisation for physics analyses [29]. A reference process can be used to estimate the cross section of a given process:

$$\sigma_{\text{exp}} = \frac{N_{\text{exp}}}{N_{\text{ref}}} \cdot \sigma_{\text{ref}}, \quad (4.6)$$

where N_{ref} and σ_{ref} are the number of reference events and the cross section of the reference process. Furthermore, the same integrated luminosity has to be used for both processes. Colliders used nowadays employ bunched beams [1]. If two bunches with n_1 and n_2 particles, respectively, collide head-on with frequency f , the instantaneous luminosity is given by:

$$\mathcal{L} = f \cdot \frac{n_1 n_2}{\sqrt{\epsilon_x \beta_x^* \epsilon_y \beta_y^*}}, \quad (4.7)$$

where $\{x, y\}$ are the coordinates transverse to the beam, $\beta_{x,y}^*$ are the amplitude functions at the interaction point, where the beam optics produces a narrow focus, and emittance ϵ is a measure of the beam width defined as $\epsilon_x \equiv \frac{\sigma_x^2}{\beta_x}$, with σ_x and σ_y being the root mean square (RMS) of the transverse beam sizes in horizontal or vertical direction, respectively. A high luminosity can be achieved with a high population of bunches of low emittance

⁴Barn b: 1 b = 10^{-24} cm^2

4. Experimental Setup

colliding at high frequency at locations where the beam optics provide low values of the amplitude functions. As the instantaneous luminosity depends on the beam parameters, it has to be measured when the beam parameter change. The instantaneous luminosity can be estimated by the event rate of a reference process if the cross section is known. The visible cross section σ_{vis} is given by:

$$\sigma_{\text{vis}} = \sigma(E) \cdot A(t, \mu, \dots), \quad (4.8)$$

where the cross sections depends on the collision energy E and the detector acceptance A depends on time t , the mean number of interactions per bunch crossing μ and other parameters. *Luminometers* are independent detectors or parts of the detector for measuring the instantaneous luminosity. Once the calibration constant σ_{vis} has been determined for a luminometer, the luminosity can be estimated using

$$\mathcal{L} = \frac{\dot{N}}{\sigma_{\text{vis}}}, \quad (4.9)$$

where $\dot{N}$ is the event rate for this luminometer. Ideally, the visible cross section of the luminometer should not be time-dependent and should not depend on experimental conditions. The visible cross section has to be determined again, when the beam parameters, including the beam energy, change. Exactly the same detector configuration has to be used during data-taking as during calibration. In summary, there are two important parts of the luminosity measurement. First, the luminometer, which measure the event rate. Secondly, the calibration of the luminometer. In the following, both parts of the luminosity measurement are introduced.

Luminometer CMS uses five detectors to monitor and measure the luminosity based on rate measurements [30]. The CMS silicon pixel detector, the DTs in the barrel of the muon system, the forward hadronic calorimeter, the Fast Beam Conditions Monitor (BCM1f), and the Pixel Luminosity Telescope (PLT) are used as luminometers. The luminometer of the pixel detector and the DTs use the standard CMS trigger and data acquisition system, while the PLT, BCM1f and HF have an independent, fast readout system. However, the silicon pixel detector and the DT have a very low occupancy and a very good stability over time [30]. In LHC Run I, only the silicon pixel detector and the HF were used as luminometer [31].

The HF has a high rate acquisition [29] and it is most sensitive to the electromagnetic component of the hadronic showers. Two methods have been studied to estimate the luminosity [29]. The first one, referred to as *zero-counting*, counts the hits above threshold of single physical towers and averages the result from each tower. The second method

exploits the linear relationship between the total transverse energy deposit in the HF and the number of interactions, consequently also the luminosity. Both methods require that the mean value of interactions μ is proportional to the luminosity.

The Pixel Cluster Counting (PCC) method employs the large number of pixels in the CMS detector. A given pixel has an exceedingly small probability to be hit by two different tracks from the same bunch crossing. For $\mu = 25$, the fraction of occupied pixels is less than a per mille [30]. It is thus expected that the number of hit pixel clusters is a linear function of the number of interactions per crossing. Thus, the number of hit pixel clusters is a good measure of the instantaneous luminosity, which is given by:

$$\mathcal{L} = \frac{\langle N_{\text{cluster}} \rangle \cdot f}{\sigma_{\text{vis}}^{\text{PCC}}}, \quad (4.10)$$

where $\langle N_{\text{cluster}} \rangle$ is the mean number of hit pixel clusters, f the orbit frequency of the LHC, and $\sigma_{\text{vis}}^{\text{PCC}}$ the visible cross section of the PCC method. Pixel modules which have not been fully operational during calibration are omitted for the offline luminosity measurement.

As the PCC method has a very small dependence on pileup and other experimental conditions, this method is chosen for the precision offline luminosity measurement. The HF measurements have a smaller statistical uncertainty and can be used for cross checks or studies of the systematic uncertainties of the luminosity measurements.

Van der Meer (VdM) scans The absolute cross section is evaluated using *Van der Meer* (VdM) scans [33], which allows to measure the luminosity per colliding bunch pair from machine parameters. A counter system measures the counting rate proportional to the rate of the beam-beam interaction. One of the two beam is displaced and the luminometer rate is measured as function of the beam-beam separation resulting in a maximum at zero separation. The VdM scan method is used under the assumption that the two bunch densities factorise in x and y [30]. These scans are performed with a dedicated LHC machine set up. Two beams are scanned through one another in the transverse plane of the detector. The luminosity and the luminosity rate must be determined at the same time to measure the visible cross section:

$$\sigma_{\text{vis}} = 2\pi \Sigma_x \Sigma_y \langle n \rangle_0, \quad (4.11)$$

where Σ_x and Σ_y are the effective beam widths and $\langle n \rangle_0 = \frac{1}{2} (R_x + R_y)$ with normalisation rates R_x and R_y , which are the amplitudes of the fitted scan curves. Beam Imaging scans are used for studies on the beam shapes. The transverse and longitudinal interaction point centroids are also determined by the VdM scan [29].

4. Experimental Setup

Luminosity integration After the VdM scan, the visible cross section of the luminometer is known and the instantaneous luminosity can be measured. The integrated luminosity is obtained by summing the luminosities of small time intervals. The convenient minimal time interval to consider for the estimation of luminosity is the *luminosity section* (LS) defined as 2^{18} LHC orbits and a corresponding t_{LS} (23.31 s in 2012). The average number of clusters per event is computed for each LS and the luminosity for the LS is derived and multiplied by t_{LS} [31].

A bunch spacing of 50 ns of the proton-proton collisions was used in 2012, when the LHC was operating at $\sqrt{s} = 8$ TeV. The Van der Meer scans were performed in November 2012. The visible cross section was measured to be $\sigma_{\text{vis}} = 7.230 \pm 0.038$ b and the overall uncertainty on the luminosity measurement was estimated to be 2.5% (syst.) + 0.5% (stat.) [31]. In 2015, a bunch spacing of 25 ns was used for the proton-proton collisions at $\sqrt{s} = 13$ TeV. The Van der Meer Scans were performed in August 2015 measuring an overall cross section of $\sigma_{\text{vis}} = 8.992 \pm 0.007$ b and the overall uncertainty on the luminosity measurement was estimated to be 2.7% [30]. This is better than the design goal of a systematic accuracy of 5% [29].

5. Event simulation

In particle physics, accurate simulations for signal and often overwhelming backgrounds are needed for measurements and searches for new physics. The basic collision, e.g. proton-proton collisions, and decay processes in an event can be described by perturbative quantum field theory. The quantum chromodynamics (QCD) bound states cannot be described with perturbative QCD (pQCD), therefore phenomenological models are needed to describe hadronisation.

Simulated particle physics events can be generated using software libraries, called event generators, which factorise the full process of the event simulation into individual tasks. The probabilistic branching between these individual problems is done by Monte Carlo (MC) methods, a class of computational algorithms that rely on repeated random sampling to have the same average behaviour and fluctuations in simulation as in collision data. MC simulations can be used to compare collision data to theoretical prediction, therefore it is very important for data analysis that the simulations are properly implemented. Events containing beyond standard model particles can be generated to see how the event signature would look like and to compare its signature to the one of generated background processes.

General-purpose Monte Carlo (GPMC) generators, like PYTHIA [57] or HERWIG [58], provide fully exclusive simulations of high-energy collisions, but there are also event generators which are specialised on a certain aspect of the event simulation. Hard-process generators use implemented perturbative matrix elements for the scattering process and hadronic event generators simulate the initial- and final-state particle showers, hadronisation, and soft hadron-hadron physics including composition and substructure of the initial state. Specialised packages for hadron decays exists. Figure 5.1 gives an overview of different steps in MC generation for proton-proton collision events. Final-state particles generated by event generators can be fed into the detector simulation. A list of event generators can be found in Table 5.1.

5. Event simulation

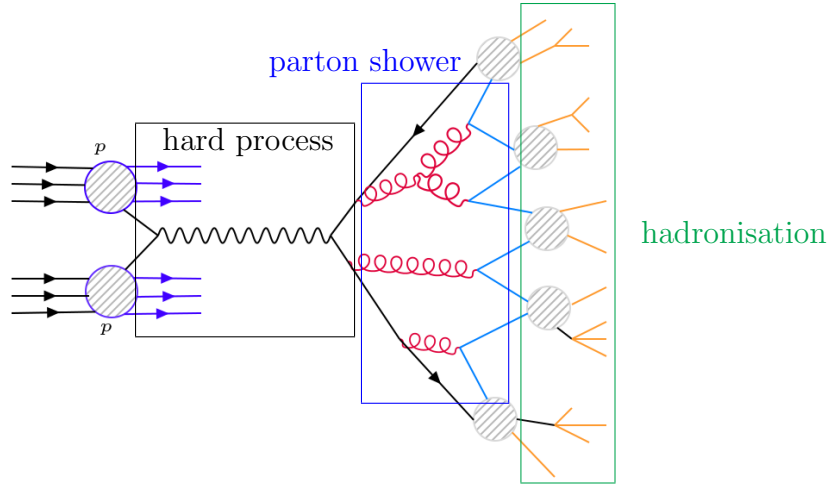


Figure 5.1.: Monte Carlo simulation of an event in proton-proton collisions

This chapter follows closely the review on Monte Carlo Event Generators [59]. First, the simulation of short distances processes is described, where the momentum scales are high enough to use perturbative quantum chromodynamics for the description of the strong interaction between quarks and gluons. Afterwards, hadronisation models are discussed which become important at lower energy scales, where quarks fragment into composite particles called hadrons. Soft hadron-hadron physics modelling, parameters and tuning are introduced. Finally, some Monte Carlo generators important for the discussed analyses are presented and the simulation of the detector response is introduced.

Table 5.1.: List of Monte Carlo generators

Generator	Purpose	Description
PYTHIA	GPMC generator	matrix element, hadronic event, particle decays
HERWIG	GPMC generator	matrix element, hadronic event, particle decays
MadGraph	matrix element	multi-purpose parton level generator (LO)
MC@NLO	matrix element	next-to-leading-order (NLO) QCD matrix elements, using HERWIG for parton shower
POWHEG	matrix element	matrix elements (QCD: NLO)
aMC@NLO	matrix element	MadGraph interface for QCD NLO calculation of matrix elements
EvtGen	decays	package for hadron decays
TAUOLA	decays	package for tau decays
GEANT4	detector simulation	detector and detector response

5.1. Short distance simulation

Short-distance components are responsible for the primary process, decay of short-lived particles and QCD & quantum electrodynamics (QED) radiation at time scales much below $1/\Lambda$, where Λ is a typical hadronic scale of a few hundred MeV, corresponding to an inverse femtometre. For momentum scales much larger than this scale, soft- and collinear-safe inclusive observables, such as total decay widths or inclusive cross sections, can be computed with pQCD theory. Large logarithmically divergent corrections, arising from final-state collinear splittings and soft emissions, cancel against virtual corrections in the total cross section. For collisions including hadrons, initial-state collinear singularities are factorised into parton density functions (PDFs), therefore the cross section for the basic process remains accurate up to higher order corrections if it is interpreted as inclusive cross section, otherwise singularities in QCD can lead to a non-convergence of the fixed-order expansion.

Hard-process generators or matrix-element generators generate the matrix element of the hard process and large-angle emissions. Hadronic event generators use collinear (small-angle) and soft (small energy) approximations to generate the parton showers. The Les Houches Event File (LHEF) standard [60] is used to transfer parton-level events from a hard-process generator to a shower generator, containing a list of particles and the used free parameters.

Matrix elements Hard-process generators generate the exact matrix elements for the production of the basic process and a certain number of n additional partons for hard, large-angle emissions. At tree-level accuracy of the hard process, radiation of extra partons is not included. Next-to-leading order (NLO) corrections can be added by including the radiation of an extra parton with tree-level accuracy and including all NLO virtual corrections. Final-state partons of the hard process and their phase space are needed by the shower algorithms as input.

Shower algorithm The cross section for a generic hard process is computed by a shower algorithm. It begins with the generation of the kinematics of the basic process, followed by a sequence of independent shower splittings. For each primary process parton i , the probability for undergoing a branching before the infrared cut-off is estimated using Sudakov form factors $\Delta_i(t, t')$ [61]. The infrared cut-off is defined by the shower hadronisation scale Λ , or the decay width for an unstable particle. The Sudakov form factor is interpreted as the probability for a splitting not to occur between two scales $t < t'$. Either a splitting into two partons is generated or the parton is defined as final parton. The probability of the parent parton i , with energy fraction z , to decay into two partons j and k is given by the Altarelli-Parisi splitting kernels [62]. This procedure is applied recursively to all

5. Event simulation

generated partons and the algorithm stops, when no final-state parton undergoes further splitting.

One can assign to each splitting vertex the energy fractions z of the two partons, the azimuthal angle ϕ of the splitting process with respect to the incoming parton momentum, and an ordering variable for the purpose of ordering the parton splittings. Different ordering variables exist. HERWIG uses the angle of emission as ordering variable and PYTHIA uses the imparted transverse momentum $p_{\perp}$. A probability is assigned to each splitting vertex to calculate the cross section for the given final state.

There are two sources of infrared singularities in massless field theories like QCD. One coming from collinear singularities, when a final-state particle becomes either collinear to another final-state particle or to an initial-state one. The other source is coming from emissions of soft gluons at arbitrary angles. Both effects can be correctly described by using the angle of emissions as ordering variable. This is done by HERWIG. Angular-ordered showers use colour information, which is tracked by MC generators. The angle between colour-connected partons determines the initial angle for the shower development. PYTHIA uses a virtuality-ordered ($p_{\perp}$ -ordered) shower evolution and an angular veto to avoid the particle multiplicity growing too rapidly with energy.

Quark masses larger than Λ , like c , b , or t quarks, act as a cut-off for collinear radiation. The divergent behaviour is regulated for angles between both produced partons $\theta \leq \theta_0 = m_q/E$, for a quark with mass m_q and energy E . Less collinear activity is expected for heavy quarks than for light ones, therefore they carry a larger fraction of the momentum acquired in the hard process. PYTHIA includes mass effects by using a matrix-element correction method, whereas HERWIG includes these effects via Altarelli-Parisi splitting kernels.

PYTHIA does not take spin correlations into account generating parton showers, while HERWIG takes them into account. Spin information can be used in analysis to distinguish between signal and other background processes, when particles with different spins are involved. Spin correlations coming from gluon splittings ($g \rightarrow q\bar{q}$, $g \rightarrow gg$) are partially included in PYTHIA by averaging over polarisations to compute the correlation between two successive splittings. Causal correlations are argued to be small and discarded. HERWIG fully includes spin correlation effects.

An incoming particle can radiate gluons before entering the hard-scattering process, therefore initial-state radiation (ISR) has to be taken into account leading to a non-vanishing transverse momentum of the particles in the matrix element. Backwards-evolution algorithms are used for considering the highest virtualities first. The main difference between algorithms are the details how the recoils are distributed among the other partons.

Shower algorithms also add electromagnetic corrections coming from photon emission from light charged particles. Soft photon emission is especially important for electrons, so that their inclusion in the simulation is mandatory. In practice, there has to be a cut-off for the electromagnetic shower in order to terminate the algorithm. For electrons, this cut-off is the electron mass and for quarks, the photon wavelength has to be smaller than the typical hadronic size of about 1 fm. Hadron and tau decays involving charged particles can produce additional bremsstrahlung.

Another important part is the decay of short-lived particles. Particle production and their decay are treated as being factorisable. The mass distribution and the distributions of its decay products are relevant. Different parametrisations can be used for the mass distribution, like a δ function at pole mass m_0 or a Breit-Wigner distribution with particle width Γ . The decay products can be described by the pole mass and a uniform phase-space distribution, or by using differential decay matrix elements. Most parton-shower models preserve the total invariant mass of the decay products to keep the original resonance shape. Decay tables are used as parametrisation for particle decays.

Matching Quarks and gluons carry a colour charge and cannot exist isolated due to QCD confinement. Jets are collimated bunches of hadrons or their decay products arising from hadronisation of a quark or a gluon. A $(N + 1)$ -jet event can be obtained by collinear/soft-radiation of an appropriate $(N + 1)$ -parton final state, generated by a matrix element generator, or from an N -parton final-state with hard, large-angle emission during shower evolution. If different generators have been used for generating matrix-elements and parton showers or extra partons have been generated by the hard-process generator, a matching has to be done between both generators.

The hard-process generator generates the exact matrix elements for the production of the basic process and up to n partons. They are tree-level accurate at large angles and match the results of the shower algorithms at small angle. Each produced parton pair has a relative transverse momentum above a scale Q_{cut} . There are no emissions below the scale Q_{cut} and no final states with more partons.

The hadronic event generator generates splittings for scales below Q_{cut} and events with less than these n additional partons. It also generates splittings for initial events with n partons for relative transverse momenta below the scale of smallest pair momentum Q_{cut} . The matching parameter Q_{cut} has to be large enough for fixed-order perturbation theory to hold, but also it has to be small enough, so that the shower algorithm is accurate for emissions below it. Consistent electroweak coupling constant, α_s , choices between real (matrix-element driven) and virtual (parton-shower driven) corrections have to be used.

Matching schemes [63] define which of both paths should be generated for a given event to avoid phase-space double-counting as well as unpopulated phase-space regions. The

5. Event simulation

choice of the path is optimised to use the best possible approximation to given kinematics. First, a jet measure is defined and all cross sections are calculated for the production of final state X and n jets. Then the hard parton samples are generated with a probability according to the total cross section and the matrix element. The events are accepted or rejected with a dynamical, kinematics-dependent probability. Finally, the parton shower is invoked. No extra jets are produced by the parton shower. The last two steps can be done in one step.

There are different matching schemes which differ how the calculation is distributed between matrix element and shower evolution. They differ in their jet definitions, acceptance/rejection of jet configurations, the starting conditions of the parton shower and the jet vetoing inside the parton shower.

5.2. Hadronisation models

The parton shower ends at the infrared cut-off Λ which is by construction identical to the hadronisation scale Q_{had} . Hadronisation describes the formation of hadrons out of quarks and gluons. GPMCs transform coloured partons into a set of colourless hadrons. This takes place at scales at long distances and low momentum transfers, where non-perturbative effects become important. There is no exact theory for hadronisation. GPMCs use QCD-inspired models which rely on the colour-flow information between partons as starting point for hadronisation. PYTHIA and HERWIG use different models for hadronisation. Fragmentation functions give the probability distribution for the hadron energy fraction from coloured partons. Hadrons can decay further, therefore the decay modelling is an important part.

Fragmentation function Fragmentation functions are probability distributions for inclusive hadron spectra. They can be defined at an arbitrary perturbative scale Q and can be calculated using pQCD and a non-perturbative initial condition obtained by fitting hadron spectra. The MC modelling includes more information and is exclusive, while the fragmentation function only describes inclusive spectra. One example is the Peterson formula [64]:

$$f(z) \propto \frac{1}{z} \left(1 - \frac{1}{z} - \frac{\varepsilon_q}{1-z} \right) \quad (5.1)$$

with z the fraction of the quark momentum which is carried by the hadron and ε_q a free parameter expected to be proportional to $1/m_q^2$.

String model String models are based on "linear confinement". In quenched lattice QCD, the potential of the colour-dipole field between a colour charge and an anticharge appears to grow linearly with charge separation at distances greater than a femtometre. One of these models is the Lund model which is implemented in PYTHIA. Colour-connected quark-antiquark pairs emerge from parton showering. When the pair moves apart, the potential between both charges grows linearly, being described by a colour flux tube being stretched between both quarks. As the string grows, the potential increases and non-perturbative creations of new quark-antiquark pairs are favoured, which break the string in two new strings. Intermediate gluons can also be involved, representing them as kinks, building a transverse structure in the one-dimensional object, while infinitely soft gluons are absorbed into the string.

String breaks are causally disconnected. In the Lund model, string breaks are generated starting with the leading hadrons, containing the endpoint quarks, iterating inwards to the centre of the string, alternating between both sides. A single on-shell hadron can be split off in each step. Quarks can acquire a transverse momentum $p_\perp$ by the breakup process. The production of quarks is exponentially suppressed with m_q and $p_\perp$. Charm and bottom quarks are not produced in soft fragmentation due to their large mass. Value of strangeness suppression is determined from experimental observables using ratios of strange to non-strange hadrons. A small amount of strangeness is generated in the parton shower through perturbative $g \rightarrow s\bar{s}$ splittings. Baryons can be produced by allowing string breaks to of pairs of diquarks, which are loosely bound states of two quarks. The relative rate of diquark to quark production can be extracted from collision measurements of the proton to pion ratio p/π .

In the next step, the produced quarks are assigned to hadron multiplets, mesons and baryons. Individual rates are not predicted by the model leading to many free parameters. Another parameter is the ratio of vector to pseudoscalar production V/P . In the final step, the fraction z of the quark longitudinal-momentum carried by the created hadron is estimated using the Lund symmetric fragmentation function:

$$f(z) \propto \frac{1}{z}(1-z)^a \exp\left(-\frac{b(m_h^2 + p_{\perp h}^2)}{z}\right) \quad (5.2)$$

which is normalised to unit integral. The hadron mass m_h and the transverse momentum $p_{\perp h}$ are fixed parameters. The dimensionless parameter a can be flavour-dependent and dampens the hard tail of the fragmentation function, while the parameter b has the dimension GeV^{-2} and is a universal constant related to the string tension. Heavier flavoured hadrons carry a larger momentum fraction z of the heavy quark [65], therefore their fragmentation functions should be harder than the one of a light hadron.

5. Event simulation

Cluster model Cluster models are based on pre-confinement [66], i.e. the colour structure of a pQCD shower evolution at any scale Q_0 can be described by colourless clusters, which are formed from colour-connected quarks. There are meson-like clusters and baryon-like clusters. Their invariant mass distribution only depends on Q_0 and the fundamental QCD scale Λ_{QCD} , and not on the hard subprocess scale Q at which the quarks and gluons are produced, for $Q \gg Q_0 \gg \Lambda_{\text{QCD}}$. HERWIG implements a version of the cluster model.

At the end of the parton shower, all gluons are forced to split into quark-antiquark pairs. A new set of colour-singlet clusters with low-masses is obtained, formed only by quark-antiquark pairs. Flavour and momenta are assigned to the produced quark pairs. Clusters with an invariant mass above the cut-off value, typically 3-4 GeV, undergo sequential $1 \rightarrow 2$ cluster breakups, along an axis defined by the constituent partons, until all sub-cluster masses are below the cut-off value. Some clusters are allowed to decay directly to a single hadron and the nearby clusters absorb any excess momentum. The probability for single-hadron cluster decays is an additional parameter. Remaining clusters are interpreted as smoothed-out spectrum of excited mesons, each of which decays isotropically to two hadrons. Relative probabilities are proportional to the available phase space, and no further parameters beyond the spectrum of hadron masses have to be introduced. Baryons can be produced by allowing clusters made of diquark and anti-diquark pairs.

Hadron and tau decays Many primary hadrons, originating from string breaks or cluster decays, are unstable and decay further, until a set of particles is obtained which is stable on relevant time scales. The decay modelling can have a significant impact on final particle yields and spectra. Particle summary tables represent a summary of available experimental measurements and may be incomplete or inconsistent within experimental precision. An example are charmed hadrons, hadrons which contain a charm quark. Many charmed hadron decays have been measured at different experiments. All measured branching ratios do not have to sum to unity, but MC simulations need decay packages with quantified, consistent information, with all branching fractions adding up to unity. Thus, choices have to be made when adapting summary tables and double counting has to be avoided. Another decision has to be taken for the differential distribution of the decay products in phase space. Matrix elements can be used, but this can be done only for a selected class of decays. Generated events can be reweighted using the obtained particle momenta, or external packages, such as EvtGen [67] for hadron decays or TAUOLA [68, 69] for tau decays. It depends on the generator if additional effects, such as B-meson oscillations, or CP-violating effects, are included.

5.3. Soft hadron-hadron physics modelling

The basic process, its parton-shower evolution and hadronisation are not the complete picture of a hadron-hadron collision. Additional activity beyond the basic process and its associated ISR and FSR activity is referred to as underlying event (UE). It is expected that the dominant part is coming from additional colour exchanges between the beam remnants. Multiple parton-parton interactions (MPI) can produce two or more back-to-back jet pairs with each pair having a small transverse momentum, while jets coming from bremsstrahlung tend to be aligned with their parent parton. Most MPI are soft. They influence the colour flow and the total scattered energy of the event. This affects the final-state activity and increases the particle multiplicity in the final state.

Hard jets appear to sit on top of a higher "pedestal" of underlying activity than events with no hard jets. This is called the "jet pedestal" effect and comes from the impact-parameter-dependence, since central collisions are more likely to contain at least one hard scattering due to the higher probability of interactions. The impact-parameter shape is therefore another free parameter.

5.4. Parameters and tuning

Very important for event simulation is the accuracy of the used models, which depends on the inclusiveness of the chosen observables and on the sophistication of the simulation. Improvements of the theoretical calculations can improve the models, but the precision also depends on the constraints in the free parameters. Existing collision data is used to constrain them. This is referred to as generator tuning. It makes no sense to tune a MC generator beyond the constraints in theoretical and experimental precision due to overfitting, otherwise it would describe fluctuations or noise instead of the underlying relationship. Parameters are α_s , PDFs for the evolution probability for ISR, properties of non-perturbative fragmentation functions, the matching parameter Q_{cut} and the hadronisation scale Q_{had} . Other parameters are the production rates of mesons and baryons, including the ratio of vector to pseudoscalar production, their masses and decay widths. These parameters and the parameters of the decay modelling give the largest amount of free parameters. Thus, the hadronisation parameters should be retuned if significant changes to decay treatment are made. The impact parameter shape of the beam particles is an important parameter for the underlying event activity and the jet-pedestal shape.

The event modelling and the generator tuning has an influence on the final particles and their spectra. Events generated with different generators or tunes can be different and might not describe the collision data in the entire phase space.

5.5. Monte Carlo generators

Since the full process can be factorised into individual tasks, MC generators can either simulate the full process, like the GPMC generators PYTHIA [57] and HERWIG [58], or they are specialised on a certain aspect of the event simulation. There are many MC generators on the market that use different methods. In the following, some of these MC generators are described in more detail. They have been chosen according to the relevance for this thesis.

First, matrix-element generators MadGraph [70], POWHEG [71], and aMC@NLO [72,73] are described. MadGraph generates the matrix element with leading-order (LO) accuracy and POWHEG with next-to-leading order (NLO) accuracy. MadSpin [74], which is a MadGraph tool for resonance decays which preserve spin correlations, is also described. Then the MLM matching scheme [63, 75, 76] and the FxFx merging scheme [77] are presented. PYTHIA [57] and HERWIG [58] have been introduced before, but a more detailed summary is now given. Finally, two MC generators specialised on simulation of particle decays are introduced: EvtGen [67], a MC generator for simulation of B-meson decays and decays of other resonances, like charmed hadrons, and TAUOLA [68,69] which is specialised on tau decays.

MadGraph 5 is a matrix element generator for high energy physics processes [70]. These processes can involve final states with a large number of jets, heavy flavour quarks, leptons and missing energy. Events can be generated for any new physics model which is renormalisable or an effective field theory and can be written in form of a Lagrangian.

All local field theories are defined by a set of particles with certain quantum numbers, a set of parameters, and a set of interactions among these particles. The model can be implemented by using FeynRules [78], which derives the Feynman rules from a given input Lagrangian. Model restriction files can be used to fix some parameters at some given value, which is often needed in phenomenological applications.

A process is defined in terms of initial- and final-state particles. MadGraph generates the Feynman diagrams and computes the code which is necessary to evaluate the matrix element at a given phase-space point for a process using the Feynman rules. There are no restrictions on the number of particles in an interaction vertex. This is needed for higher-dimension effective interactions with five or more fields. The diagram generation algorithm only constructs topologies that will yield a valid diagram and recursively generates all diagrams in parallel reusing already calculated subdiagrams. Core processes can have a number of decay processes. Generation of decay chains is done successively. Helicity amplitudes are generated separately for each process. Helicity wave-functions are generated for each external leg in a diagram. Wave functions corresponding to identical subdiagrams,

but contributing to different diagrams, can be reused. A fermion flow is defined for each fermion line in a diagram. Fermions are grouped in pairs with each constituting a fermion line. Colour coefficients in scattering amplitudes of states which are charged under $SU(3)_C$ are computed. Multiparton amplitudes, i.e. amplitudes with many quarks or gluons, are important as they correspond to leading order approximation of multi-jet production. Exotic colour representations or non-standard colour structures in interaction vertices can be implemented, because of the modular structure of MadGraph. Colour objects and their algebra are coded in a symbolic way and the full amplitude is split into gauge invariant subamplitudes. The colour matrix is calculated and stored in the output. The resulting matrix element contains the full spin correlation and Breit-Wigner effects, but are not valid far from the mass peak.

The POWHEG BOX is a framework for implementing next-to-leading order (NLO) matrix-element calculations [71]. It includes radiation of an extra parton in the matrix element and NLO virtual corrections. The generated events have a positive weight. POWHEG can be interfaced with parton shower MC programs, such as PYTHIA or HERWIG, and it does not depend on the MC program used for showering. The POWHEG method has been applied to several hard processes, including single top production, $t\bar{t}$ production and Higgs production via gluon fusion, vector boson fusion, and Higgs boson fusion production associated with a vector boson (Higgs strahlung).

POWHEG needs the LO matrix-elements, the real matrix-element and the finite part of the virtual corrections as input. With this information, POWHEG finds all singular regions. The real cross section is split into a sum of contributions that are divergent in one singular region only using the Frixione, Kunszt and Signer (FKS) subtraction scheme [79, 80]. In the FKS framework, singular regions are characterised by a final-state parton becoming collinear or soft to either an initial-state parton or a final-state parton. The underlying LO-diagram is obtained by replacing this parton pair with a single parton of appropriate flavour, so that singular regions can be grouped according to their underlying LO-diagram. Only emitted final-state partons can lead to soft singularities.

After finding all singular regions, soft and collinear counterterms as well as remnants are built. Radiation is generated using the POWHEG Sudakov form factor. Initial-state radiation (ISR) and final-state radiation (FSR) are treated differently. Born-like events are generated, when no radiation is generated. Gluon splittings into heavy quark pairs below threshold are avoided. An azimuthal symmetry is assumed. Finally events are put into the Les Houches interface. Colours are assigned to incoming and outgoing partons on basis of the underlying LO-process and to the real emitter and the radiated parton. In case of a gluon splitting into two gluons, there are two possible colour assignments. Information on intermediate resonances is made available to the shower program, since the resonance mass

5. Event simulation

must be preserved by the shower. Also its decay products are specified. Colourless and massive coloured particles remain the same at LO and NLO level, the flavour structure has one more light parton in its final state. Virtual terms, the interference of LO and one-loop amplitude, have the same flavour structure as the LO term.

aMC@NLO automates all aspects of NLO computation and its matching with parton showers [72, 73]. NLO calculations can be achieved by computation of one-loop matrix elements, tree-level matrix elements, and combining both matrix elements. One-loop amplitudes are evaluated with MadLoop [81], and tree-level computations are performed using MadGraph. The MC@NLO formalism [82] is used for matching matrix elements and parton showers. MC counterterms are used to remove undesirable large short-scale effects and to set an upper bound for the hardness of each branching. Matched samples, which differ by their final-state multiplicity, can be merged using the FxFx- merging scheme.

MadSpin is a method which preserves spin correlations and finite width effects for narrow-resonance decays [74]. It has been implemented as a tool in the MadGraph5 framework. Spin correlations of unstable particles, which decay products are detected, are relevant for analyses that involve polarisation measurements. For example, tau polarisation can be measured to distinguish between Z- and SM Higgs-boson decays, or as probe of new physics [83]. Vector bosons decay into either $\tau_L^- \tau_R^+$ or $\tau_R^- \tau_L^+$, while the Higgs boson decays into $\tau_L^- \tau_L^+$ or $\tau_R^- \tau_R^+$.

With MadSpin, the amplitudes for production and decay of the unstable particles can be calculated separately, without losing spin correlations. The kinematics of decay products are determined given the kinematics of unstable particles. Finite width effects and off-shell effects are considered using a Breit-Wigner distribution for particle masses. It is suited for resonance decays, where the production of these resonances were generated with NLO accuracy [73, 74].

The MLM matching scheme is a matching algorithm, which has been implemented in the ALPEN generator [75, 76]. It matches partons from matrix element calculations to jets reconstructed after shower generation.

First, matrix-element generators generate the hard process. Parton-level events are required to have at least a minimum transverse energy E_T^{min} for partons and a separation greater than a minimum value $\Delta R_{jj} > R_{min}$ between them. These events are showered without any hard-emission veto during shower. Accepting/rejecting events is done after the parton shower. A jet cone algorithm with cone size R_{min} and E_T^{min} is applied. Starting with the hardest parton, the jet closest in (η, ϕ) is selected and both match if the distance is smaller than R_{min} . If they match, the jet is removed and matching is done with the

next parton. If an event has not a matched jet for each parton, this event is rejected. This is the case for collinear partons or soft partons, which do not lead to an independent jet or are too soft for jet reconstruction.

The FxFx merging scheme is an NLO merging procedure [73,77]. Exclusive events with J light jets can have an NLO accuracy by the computation based on matrix elements that have J and $(J + 1)$ partons. NLO mergings are more complicated than LO ones, because matrix elements are considered twice, as Born contribution for processes with J partons, and as real-emission contribution, IR subtraction terms, and the one-loop contributions to a processes with $(J - 1)$ partons. Double counting of matrix-element contributions has to be avoided. This can be achieved by parton-shower dependent MC counterterms and Sudakov reweighting, as Sudakov form factors include virtual corrections and non-emission probabilities. Contributions from the matrix element dominate at hard scales, while the parton shower dominates at soft scales. A probability function estimates which description is more accurate at intermediate scales. There is no emission larger than the matching scale generated by shower. Events are reweighted and a certain amount of events might carry negative weights, which has to be taken into account after MC simulation.

PYTHIA 6 is a general-purpose Monte Carlo (GPMC) event generator for simulating high-energy-physics events [57]. It has been developed for multiparticle production in e^+e^- , pp and ep collisions and simulation of jets. Standard model (SM) physics and beyond SM physics with special effort on physics of strong interaction can be simulated. PYTHIA can generate the hard subprocess, initial- and final-state parton showers, hadronisation, decays and the underlying event.

More than 300 hard processes have been implemented for generating the matrix elements for final-state and phase space calculation. It is optimised for processes with two partons initial-partons and one or two final-state partons in the hard process, $2 \rightarrow 1$ and $2 \rightarrow 2$ processes. Resonance decays are implemented for resonances with masses above the b-quark system, like the Z boson, the W boson, and the Higgs boson. Their partial width and branching fractions can be dynamical calculated as function of their mass. Often spin information is available for resonance decays leading to properly correlated angular correlations of the resonance decay products. In other cases, it is not available and the resonance decays isotropically.

A multijet structure is added by the parton shower. There are initial and final-state showers. Spin effects are not taken into account by the shower. Showers are ordered by their virtuality $p_{\perp}$. A backward evolution scheme is used for initial-state showering. Starting from the incoming parton of the hard interaction, the shower is traced backwards in time to find the parton which initiated the shower.

5. Event simulation

Dipole showering is used by PYTHIA 8 [84]. Gluon emission is not generated by splitting partons, but it is generated via dipole radiation. FSR are associated with dipoles which are stretched between a final-state parton and the "hole" left by an initial-state parton. Recoils are taken by the hard scattering subsystem having initial-state radiation unchanged.

Hadronisation is implemented using the Lund string model. The process is split into generating the fragmentation and the decays of these hadrons.

HERWIG 6 is a general-purpose MC event generator [58]. It includes simulation of hard lepton-lepton, lepton-hadron, and hadron-hadron scattering and soft hadron-hadron scatterings. Many SM and beyond SM hard processes have been implemented. The hard process is important for defining the phase space of the initial- and final-state showers. Heavy partons decay before hadronisation and can initiate showers before and after their decay. Initial- and final-state QCD radiation is generated by the parton shower. Colour coherence effects between initial- and final-state partons in all hard subprocesses and azimuthal correlations originating from gluon interference and polarisation are included. Soft gluon interference is taken into account via angular ordering of successive showers. A coherent branching algorithm is used for generating the parton showers. The angular variable ξ is defined for a parton splitting $i \rightarrow jk$ as product of final parton momenta divided by the product their energy: $\xi_{jk} = (p_j \cdot p_k)/(E_j E_k)$. For massless partons at small angles, this angular variable is proportional to the squared angle θ_{jk} between both final partons: $\xi_{jk} \simeq \frac{1}{2}\theta_{jk}^2$. Each angle must be smaller than the angle of previous branching of the parent parton. A backward evolution algorithm is used for generating initial-state showers. A cluster model is used for hadronisation, which is based on pre-confinement.

EvtGen is a MC generator for simulation of particle decays in high-energy physics [67]. Several B-hadron physics experiments, like BaBar, Belle and CLEO, have studied B-mesons in greater detail, therefore a simulation which is able to accurately describe observed distributions has been developed. EvtGen is a tool for simulating physics of B-meson decays and decays of other resonances. Angular and time-dependent correlations of sequential B decays are included and corresponding distributions are more precisely simulated. Models of time dependent CP asymmetries in neutral B meson decays are implemented and there is a large decay table for B hadron decays.

EvtGen uses decay amplitudes instead of probabilities. Amplitudes depend on the initial and final state, but they do not depend on the phase space and the particle energies and momenta. Using amplitudes instead of probabilities, which depend on the phase space, is efficient and leads to reusable decay models. The amplitude for each node in a decay tree is used to simulate the entire decay chain. The implementation of each decay amplitude is independent of how the mother particle was generated or how the daughter particles

decay. Exceptions are handled as special case in the decay table. One exception is the decay $B \rightarrow J/\psi K^*$, where the CP of the final state depends on if the K^* decays to $K_S\pi$, $K_L\pi$, or $K^\pm\pi$.

Event selection algorithm generates decays as sequence of sub-decays to be efficient and the maximum of the decay chain does not have to be known. Spin density matrices are computed and describe one daughter particle after summing over the degrees of freedom for the other particles. Decay models just have to specify the decay amplitude for each combination of mother and daughter spin states. A factor proportional to the decay rate makes the maximum decay probability of each sub-decay independent of the full decay chain. Kinematics are generated according to phase space and probability. Depending on this probability the event is accepted or rejected.

CP violating decays and oscillations of neutral B-mesons ($B - \bar{B}$ mixing) have been implemented in EvtGen. Development of models has focused on B-meson decays at the $\Upsilon(4S)$ resonance, which is just above the threshold to decay into two B-mesons. There are still some issues with the simulation of B_s -meson and B-baryon decays, but the B_s mixing is correctly treated. EvtGen has an efficient simulation of complex decay chains for B-mesons. Accurate simulation of B-meson decays is important for identifying b-quark jets, because algorithms for b-quark jet identification make use of B-hadrons.

TAUOLA is a library of MC programs to simulate decays of polarised taus [68,69]. Since taus are heavier than electrons and muons, they have more decay modes, and therefore are difficult to handle analytically. Their decay products are sensitive to the polarisation of the tau. Precise simulation of production and decays of taus is important, since taus can be used to test electroweak theory and they are background to other processes. An important background including taus is $Z \rightarrow \tau\tau$. Simulation of taus are also important for the Higgs boson decay $h \rightarrow \tau\tau$.

TAUOLA simulates the tau decays and it can be attached to MC programs simulating the production of taus. The full final-state topology is taken into account, including neutrinos, resonant distributions for intermediate particles and the complete spin structure. The calculation is based on quantum mechanics (QM), which is important for including spin effects. The phase space is treated independently from the matrix element and masses of all final-state particles are taken into account in the calculation. Spin polarisation is taken into account for primary and secondary decay products. A spin density formalism is used. Density matrices describe mixed states of quantum systems. This spin density formalism is analogous to probability distributions of position and momentum (phase-space) in classical statistical mechanics. More than 20 decay channels are implemented, including lepton and several hadronic modes, one-meson modes, two-meson modes, three-meson modes, 4-pion modes, and multi-pion modes. Decays with more than four taus are treated in a

5. Event simulation

simplified way, neglecting spin distributions, but including effects due to neutrino masses. The spin polarisation can be transmitted to TAUOLA as input parameter. Masses, widths and coupling constants of taus and their decay products can be adapted.

5.6. Detector simulation

In particle physics events are generated by an event generator and passed to detector simulation. Interactions of particles with the detector material and the detector response are simulated. Afterwards, these events can be reconstructed and analysed. Detector simulation is done using Geant4 [85], which is a toolkit for simulating passage of particles through matter and for simulating particle interaction in matter. Several physics models are implemented to handle interactions of particles with matter across a very wide energy range. It is used in particle physics, nuclear physics, accelerator design, space engineering, and medical physics, where large-scale, complex simulations are needed.

Detector geometry and materials are defined by the user. The geometrical model can have a large number of components with different shape and materials. Sensitive elements can be defined which record information, *hits*, needed to simulate the detector responses, *digitisation*. The geometrical structure of the detector is divided into logical and physical volumes. Logical volumes have a certain shape and they can hold other logical volumes inside. They contain the information of the material and the sensitive detector behaviour. The material can be a single element or a mixture of different elements and isotopes. Physical volumes carry information about the spatial positioning or placement of the logical volume with respect to an enclosing mother logical volume.

Particles can decay while they are transported through the geometry or they can interact with the detector material. Cross sections can be derived by formulas, parametrisations or databases can be used. Final states are separated from cross-section calculation. Both depend on energy, range, particle type and material. There are many processes with a particular initial and final state with well defined cross sections. A model can be implemented by different processes in Geant4 depending on the energy or particle type. There are electromagnetic, hadronic and optional processes. Electromagnetic processes describe electromagnetic interactions of leptons, photons, and hadrons. Geant4 can handle ionisation described by energy loss and range tables, bremsstrahlung, pair production of electron-positrons from muons, photo-electric effect, pair conversion, annihilation, synchrotron and transition radiation, scintillation, refraction, reflection, absorption, the Cherenkov effect, and many other processes. Hadronic shower can be simulated with their inclusive cross sections and final states. Particles with their basic properties, like

mass, charge, list of sensitive processes, can be defined. Particles can decay and the branching fractions for its decay products are given by theory, parameters or data-driven distributions.

Particles are transported in steps and they are tracked through materials and external electromagnetic fields. All physics processes associated with the particle can propose a step. If the last step is reached the particle is not transported further, otherwise steps proposed by actions of all attached processes are executed. Particles can decay at rest. They can loose energy and produce secondary particles, like continuous Cherenkov radiation, along the step. Some actions are invoked at the end of the step, like secondary particle production due to a decay or an interaction of the particle. When a particle or track passes a sensitive detector, hits are generated to record information needed to simulate the detector responses. Hits are a snapshot of a physical interaction or interactions of a track or tracks in a sensitive detector component. They are collected in an event object at tracking time. Digits are created from one or more hits or other digits to simulate the detector output. Digitisation is invoked by users code and can be postponed for a better pileup simulation due to overlapping signals.

Event data is generated during simulation. First, events contain primary vertices and primary particles before processing an event. After processing, hits and digitisations generated by simulation are added. Trajectories of simulated particles can be added optionally for recording of "simulation truth".

6. Event reconstruction

This chapter describes how particles are reconstructed using information from different subdetectors. A precise and reliable reconstruction over a large energy range and at high luminosities is necessary for precision measurements and to distinguish SM signatures from signatures of other new physics.

Figure 6.1 shows the event signatures of particles in the CMS detector. Trajectories of charged particles can be detected with the silicon-tracking detector and reconstructed as *tracks*. Such particles are electrons, muons, and charged hadrons. Production vertices of these particles can be also reconstructed using the tracking detector. Secondary vertices originating from bottom hadrons are especially important for identification of b-quark jets, which employs the long life time of bottom hadrons leading to a vertex, which can be reconstructed and resolved with respect to the vertex of the primary interaction. Electrons interact with the silicon tracker and the ECAL, while photons only interact with the ECAL. Other particles which only interact with the calorimeters are neutral hadrons, which primarily interact with the HCAL. On the other hand, charged hadrons do interact with the tracking detector, the ECAL and the HCAL. Hadrons originate from the hadronisation process of quarks. These hadrons can be reconstructed as *jets* also including their decay products, electrons, muons and photons. Muons traverse the whole CMS detector and their trajectories can be reconstructed in the tracker and in the muon system, the most outer part of the detector. They hardly leave energy in the calorimeters. Neutrinos do weakly interact only, therefore they cannot be directly reconstructed. However, the sum of missing transverse energy in an event can be reconstructed as a reliable signature of neutrinos or similar weakly interacting particles.

First, the reconstruction of particle trajectories as tracks is discussed and it is described how primary vertices, originating from particle decays, are reconstructed [86]. Then, electron reconstruction and identification [54, 55, 87] is discussed, followed by a discussion on the reconstruction and identification of muons [88]. Both algorithms are important parts of the search discussed in this thesis. The tag-and-probe technique [89] is introduced, and it is described how reconstruction, identification, and trigger efficiencies of leptons can be estimated using this technique [54, 55, 87, 88]. This technique is employed for

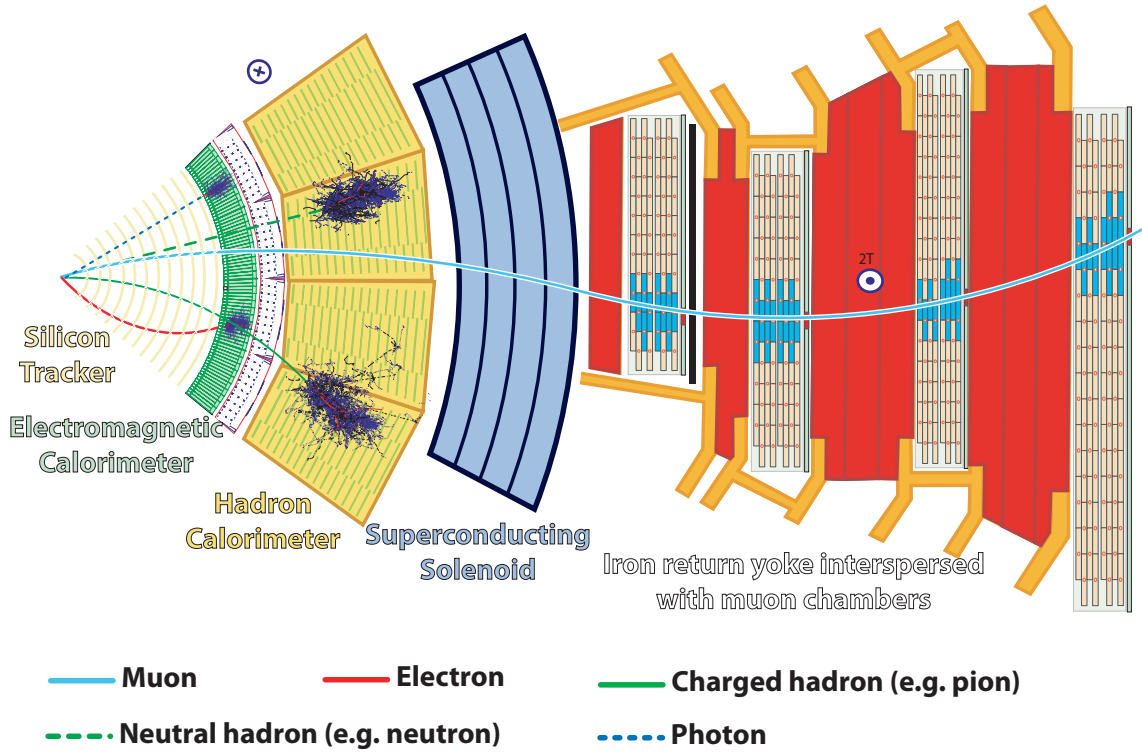


Figure 6.1.: Signatures of different particles in the CMS detector [90]. Particles which interact with the silicon tracker are shown as full line. Electrons are indicated as red line, muons as blue line, and charged hadrons as green line. Other particles, which only interact in the calorimeters, are shown as dashed lines, such as neutral hadrons (green dashed-line) and photons (blue dashed-line).

the estimation of the trigger efficiency of electrons used in the search at $\sqrt{s} = 13$ TeV. Particle-flow event reconstruction is introduced as it improves the reconstruction and identification of particles, especially of jets, by combining information from different subdetectors [91]. Then, jet reconstruction and calibration [92] are discussed. Jets are the experimental signatures of energetic quarks and gluons, which are not directly observable but fragment into a collimated bunch of hadrons flying roughly in the same direction. Among such hadrons are hadrons with a charm or bottom quark content. Such hadrons have long life-times and the vertices of their decay can be reconstructed with the CMS detector. Particles which only interact via the weak interaction, such as neutrinos, do not produce a direct response in the CMS detector, but they lead to an imbalance of the total transverse momentum, which is reconstructed as missing transverse energy [93], which is also discussed in this chapter. Finally, the identification of b-quark jets which contain bottom hadrons is discussed [94–96].

6.1. Track and primary-vertex reconstruction

Charged particles have bent trajectories in the CMS magnetic field and can be detected by the pixel and strip silicon detectors. When a particle trajectory crosses a track layer, *hits* can be reconstructed with space and time information. Combining these hits to *tracks* gives the reconstructed trajectory of the charged particle.

Multiple proton-proton interactions (*pileup*) lead to high track multiplicities, so that combinations of unrelated hits becomes more probable. The fraction of those falsely reconstructed tracks (*fake tracks*) has to be kept low by the track reconstruction algorithm, while keeping the track-finding efficiency high. The algorithm has also to be fast, since there are many tracks in an event and those tracks are needed by the other parts of the event reconstruction. Searches for high-mass dilepton resonances require a good momentum resolution up to 1 TeV. On the other hand, efficient reconstruction of low-momentum tracks, with momenta as low as 100 MeV, is needed for reconstructing hadron trajectories in order to achieve a good jet resolution. The reconstruction of 3-prong tau-decays requires a good resolution of the nearby tracks. Identification of jets originating from b-quarks rely on a very well determined impact parameter resolution of tracks from bottom-hadron decays with respect to the primary proton-proton interaction vertex.

Hits are reconstructed in the pixel silicon and strip silicon detectors. Then, trajectories of charged particles are reconstructed by the CMS tracking software, the *Combinatorial Track Finder* (CTF) [86]. The CTF is an adaption of the combinatorial Kalman filter [97–99], which is an extension of the standard Kalman filter [100] to allow pattern recognition and track fitting to occur in the same framework.

6.1.1. Hit reconstruction

A charged particle traversing a tracker layer leaves a signal in one or more pixel sensors or strip readout-chips. If the signal is above a specified threshold, these signals are clustered into *hits* and their position and uncertainty is estimated using a local orthogonal coordinate system (u, v) in the plane of each sensor.

Pixel-clusters are formed from side-by-side and corner-by-corner adjacent pixels. There are two algorithms for estimating the pixel-cluster position. A fast one used for track seeding, presented in Section 6.1.2, and pattern recognition, and a more precise one for final track fitting. Strip channels are used for hit reconstruction if their charge exceeds five times the channel noise or if its charge and the one of its neighbour both exceeding a smaller charge threshold. Clusters are seeded by channels, passing a certain charge threshold, and neighbouring channels are added if their signal is above another noise threshold. Only clusters are kept, whose charge exceeds a charge threshold. The hit position is the charged-weighted average of all strip positions in the cluster.

6. Event reconstruction

The *hit efficiency* is the probability to find a cluster in a given silicon sensor that has been traversed by a charged particle. This efficiency is determined for the pixel and strip detector using specific track collections. In the pixel silicon detector, this efficiency is estimated using tracks which either have a hit associated in the layer under study or a hit within $500\text{ }\mu\text{m}$ of the predicted track position. In the strip silicon detector, tracks which have a minimum of hits in the pixel and strip detectors are used. The hit efficiency depends on the instantaneous luminosity and the trigger rate and is on average larger than 99%. The *hit resolution* can be estimated using residuals defined by the difference between measured and expected hit-position as predicted by the fitted track. Each track is refitted excluding the hit under study. The uncertainty in the track position in the strip detector is larger than the inherent resolution due to multiple scattering. The sensor resolution depends on the cluster width and on the pseudorapidity, since the energy deposit depends on the angle at which the charged particle crosses the sensor plane.

6.1.2. Track reconstruction

Reconstructed hits are used for track reconstruction and to obtain the momentum and position parameters of the charged particle. A translation between the local coordinate system of the hits and the global coordinate system of the track is done.

A collection of reconstructed tracks is produced by *iterative tracking* with the *Combinatorial Track Finder* (CTF) reconstruction sequence [86]. Hits associated with high-purity tracks found in previous iterations are removed to reduce the combinatorial complexity for searching for more difficult classes of tracks, such as low-momentum tracks or displaced tracks. Track seeds are generated to provide initial track candidates. Then, a search for additional hits is done by extrapolating the trajectory along the expected flight path in the track finding step. The trajectory is fitted to the hits in the track-fitting step. Finally, quality selection criteria are applied and the track is either saved or discarded.

Seed generation Seed generation provides initial track candidates for an initial estimation of the trajectory, its parameters and its uncertainties. Only a few hits are used. Charged particles in an uniform magnetic field follow helical paths, therefore five parameters are needed to define their trajectory. Three 3-D hits, or two 3-D hits and a constraint of the trajectory origin are sufficient for track initialisation.

Information on the position of the reconstructed beam spot and the location of the primary vertices are needed for seed generation. Seeds are constructed in the inner part of the tracker and the track candidate is built outwards. The track density is higher in the inner part, but the high granularity of the pixel detector leads to a lower channel occupancy. Other reasons for doing the seeding in the inner part of the tracker instead of in the outer part are the significant energy loss of electrons due to bremsstrahlung, and the deflection of particles with low momenta due to the strong magnetic field.

Track finding Track finding is based on the Kalman filter method [97–100]. Seed trajectories are extrapolated along the expected flight part to search for additional hits that can be assigned to the track candidate. Candidates are build by adding hits from successive detector layers and updating the track parameters at each layer making use of the location of the detected hits, their uncertainty, and the amount of crossed material at each layer.

Taking into account the current track uncertainty, it is evaluated which adjacent layer intersects with the extrapolated trajectory. Then, a search for compatible silicon modules is performed. Groups of hits are formed and the Lorentz drift of ionisation-charge carrier is recalculated using the trajectory direction on the sensor surface. A χ^2 -test is used to check which hits are compatible with the extrapolated trajectory. The trajectory is updated and new track candidates are built from each candidate by adding exactly one of the compatible hits. Track candidates found at each layer are propagated to the next compatible layers and this procedure is repeated until a termination condition is satisfied, such as the end of the tracker, too many missing hits, or too low track momentum. When the search to the outward direction gives a minimum number of valid hits, an inward search is done. Additional hits can be found, for example hits in layers closer to the interaction region than the seeding layer. The signal-to-background ratio, the impact-parameter resolution and the momentum resolution are improved through this procedure. After all tracks in a given iteration are found, a trajectory cleaner is applied to remove tracks, which share to many hits with another track candidate. Finally, requirements on the number of missing hits, on the number of hits, and the track momentum are applied. A collection of hits and an estimate of track parameters for each trajectory are the output of the track finding step.

Track fitting All trajectories are refitted for the best estimate of their parameters. The full information without the bias coming from the used beam-spot position in the seeding stage is given as input for a Kalman filter and smoother.

The Kalman filter is initialised at the location of the innermost hit. Then, the trajectory is extrapolated from one hit to the next one iteratively updating the trajectory estimate by proceeding through the full list of hits starting from the inside outwards. Material effects and inhomogeneities of the magnetic field are taken into account. For each valid hit, the hit position uncertainty is re-evaluated using current values of the track parameters. In the smoothing stage, a second filter is initialised with the results from the first iteration. The fit runs backward towards the beam-line. The weighted average of track parameters of both filters gives the optimal track parameters at any point, including the innermost and the outermost hit.

After the filtering and smoothing steps, a search for incorrectly assigned hits, outliers, is done. Those hits are removed and the track is fitted and smoothed again. This continues until no more outliers are found.

Track selection A track selection is applied to remove fake tracks. Tracks that fail certain criteria are discarded from the track collection. Such criteria are the number of layers which have a hit assigned to the track, the χ^2/dof of the fit, or the compatibility of the track to originate from a primary vertex. The fraction of fake tracks is exponentially reduced as function of the number of layers in which a track has associated hits. After the last iteration of track reconstruction and applying the track selection for each iteration, tracks found by each iteration are merged into a single collection.

6.1.3. Special track reconstruction

Electrons lose energy primarily through bremsstrahlung rather than ionisation. This leads to large energy losses before the electrons reach the ECAL - one third of the electrons radiate more than 70% of their initial energy leading to a non-Gaussian energy loss distribution. Since Kalman filter work optimal for variables with Gaussian uncertainties, the efficiency and resolution of standard tracking is not good enough for electrons. Bremsstrahlung photons emitted by electrons strike the ECAL at very similar η value, but different ϕ values. A dedicated track reconstruction for electrons is described in Section 6.2.1.

The HLT has to be able to process large amount of data in a short time. Tracking information can be used to identify interesting events, e.g. events with isolated leptons when tracks are counted within a cone around the lepton. A track reconstruction similar to the offline one is used, but more than ten times faster. Using only pixel hits, regional track reconstruction and a single iteration can speed up the algorithm. Additionally, track reconstruction is only performed after other requirements have been satisfied, e.g. electron HLT objects are certified.

6.1.4. Primary-vertex reconstruction

Location of all proton-proton interaction vertices are reconstructed in each event, including vertices originating from pileup collisions. Reconstructed tracks which are consistent with being produced promptly in a primary interaction are used. Then, tracks which appear to originate from the same vertex are clustered according to their z -coordinate at the point of closest approach to the centre of the beam spot. The clustering algorithm has to resolve nearby vertices, but it also has to take care that a vertex is not falsely split into two vertices [86]. Final assignment of tracks is decided in the last clustering step. Finally, a fit for the position of each vertex is done using an *adaptive vertex fitter* [101]. Vertex parameters, such as x , y , and z position and the fit parameters, are estimated. A weight is assigned to each track corresponding to the likelihood that it belongs to a certain vertex. A similar approach can be used for secondary-vertex reconstruction. The primary-vertex reconstruction efficiency defines how often a vertex is reconstructed successfully and if its position is consistent with its true value. This efficiency is close to 100% when more than two tracks are used for the fit.

A fast primary-vertex reconstruction based on pixel hits is used by the HLT. The impact parameter of pixel tracks relative to their vertex can be used to identify displaced tracks expected from bottom-hadron decays. Pixel tracks have a lower reconstruction efficiency than standard tracks. The primary-vertex resolution depends strongly on the number of tracks used in the fit and their momenta.

6.1.5. Beam spot reconstruction

The beam spot gives a 3-D profile of the luminous region, where LHC beams collide in the CMS detector. Beam spot parameters are determined by averaging over many collisions. The position of the beam-spot centre corresponds to the centre of the luminous region. Its r and z values are measured for each event. The position of the beam spot is important for estimating the position of the interaction point for primary-vertex reconstruction. It is used especially by the HLT.

There are two methods for measuring the beam spot position. The first one uses reconstructed primary-vertices and estimates the mean of x , y , z and the size of the luminous region through a fit of a likelihood to the 3-D distribution of the vertex positions. The other method makes use of the correlation between the transverse impact parameter d_0 and ϕ of a track belonging to the primary-vertex, when the centre of the beam spot is displaced relative to the expected position. An iterative χ^2 fit is used. Width and length of the luminous region cannot be obtained with the same algorithm, therefore a combination of both methods is used. The d_0 - ϕ fit measures the centre of luminous region in the transverse plane, x and y coordinates, and the slopes dx/dz and dy/dz . Primary-vertices are used to determine the z coordinate and RMS width ($\sigma_x, \sigma_y, \sigma_z$). The beam spot position is measured in every *luminosity section* (LS), which corresponds to events collected during a period of 23 seconds. Results of all LS intervals in a run are combined to extract the final beam spot values.

The size of the beam spot is determined with two methods. The first one uses the distribution of reconstructed primary-vertices and the second one uses event-by-event correlations between transverse impact parameters of two tracks originating from the same vertex. The beam spot size is independent of the resolution of the vertex position and impact parameters.

6.2. Electron reconstruction

A good electron reconstruction and identification is important in many analyses. The decay of the Higgs boson into four leptons $H \rightarrow ZZ^* \rightarrow 4l$ can have four or two electrons in the final state. One of these electrons can have a very small momentum requiring a good performance down up to $p_T \approx 5$ GeV. On the other hand, there are searches for high-mass resonances or other new processes beyond the standard model involving electrons with momenta above 100 GeV. A good performance over a large momentum range is therefore needed.

Electron reconstruction [87] combines the energy information of the electromagnetic shower initiated by the electron and the momentum measurement of the electron track in the tracking detector. One challenge is the precise measurement and reconstruction of electrons in the tracker. Its material causes significant bremsstrahlung along the electron trajectory and this bremsstrahlung is spreading over a large volume due to the magnetic field. Bremsstrahlung leads to an energy loss of the electrons and an early development of an electromagnetic shower in the tracker. Performance measurements have been done at $\sqrt{s} = 7$ TeV [54] and $\sqrt{s} = 8$ TeV [55].

In this section, the reconstruction of electrons is described first. Then, identification observables used for background suppression and rejection of electrons from photon conversion into an electron-positron pair are introduced. Finally, isolation observables which can be used to suppress electrons within jets are discussed.

6.2.1. Electron reconstruction

Electrons are reconstructed by combining the ECAL calorimeter information with tracker information of the tracking detector. Energy originating from the electron are clustered together in the ECAL. Tracks are reconstructed in the tracking detector and associated with energy clusters in the ECAL [54,55,87]. A standalone approach [87] is combined with the global *particle-flow* (PF) algorithm [91] for a better performance. The PF algorithm is discussed in Section 6.5.

Clustering of electron energy in the ECAL

Electrons can radiate bremsstrahlung photons. In the magnetic field, electron trajectories are bent, therefore the energy of the electron is spread in ϕ over several ECAL crystals. The spread in η of the radiated photons is negligible for electrons with $p_T > 5$ GeV. Effects induced by bremsstrahlung can be large depending on the amount of material along the electron trajectory, on average 33% ($\eta \approx 0$) to 86% ($|\eta| \approx 1.4$). In the clustering step,

6. Event reconstruction

η , ϕ and transverse energy E_T are defined relative to the centre of the detector. There are two standalone clustering algorithms and a PF-reconstruction algorithm.

Hybrid algorithm: This algorithm exploits the geometry of the EB and properties of the shower shape. Energy is collected in a small window in η and an extended window in ϕ . The crystal, containing most of the energy deposited in the considered region, is taken as seed crystal if its transverse energy $E_{T,\text{seed}}$ exceeds a certain threshold $E_{T,\text{seed}}^{\min}$. Then, arrays of $\eta \times \phi = 5 \times 1$ crystals are added around the seed crystal in both ϕ directions if they are in range N_{steps} and if their energy exceeds $E_{\text{array}}^{\min}$. Adjacent arrays are grouped into clusters. If these clusters have a seed array with an energy exceeding $E_{\text{seed-array}}^{\min}$, they are collected in the final global cluster, the *supercluster* (SC).

Multi-5×5 algorithm: This algorithm is used for energy clustering in the EE. The ECAL crystals with the maximal local energy relative to their four direct neighbours are taken as seed crystal if E_T exceeds $E_{T,\text{EEseed}}^{\min}$. Then, the energy is collected in clusters of 5×5 crystals, that can partly overlap, starting with the crystal with largest E_T . Within a range around a seed crystal in η of $\pm\eta^{\text{range}}$ and in ϕ of $\pm\phi^{\text{range}}$, clusters are collected into an SC if their E_T is larger than $E_{T,\text{cluster}}^{\min}$. Threshold values are listed in Table 6.1.

The preshower energies are added to the SC by extrapolating the energy-weighted positions of all clusters in the SC to the planes of the preshower. The preshower clustering range in ϕ is defined by the maximum distance in ϕ between the clusters and the most energetic cluster. This value is extended by ± 0.15 rad. A range of ± 0.15 is used in η with respect to the most energetic cluster.

Table 6.1.: Threshold values used in the clustering algorithms for proton-proton collisions at $\sqrt{s} = 8$ TeV [55].

Barrel		Endcaps	
Parameter	Value	Parameter	Value
$E_{T,\text{seed}}^{\min}$	1 GeV	$E_{T,\text{EEseed}}^{\min}$	0.18 GeV
$E_{\text{seed-array}}^{\min}$	0.35 GeV	$E_{T,\text{cluster}}^{\min}$	1 GeV
$E_{\text{array}}^{\min}$	0.1 GeV	η^{range}	0.07
N_{steps}	17 (≈ 0.3 rad)	ϕ^{range}	0.07 rad

Both clustering algorithms group the ECAL energy into an SC, with the *SC energy* E_{SC} corresponding to the energy sum of all its clusters. The energy-weighted mean of the cluster positions defines the *SC position*. Table 6.1 lists the threshold values used in the clustering algorithms for proton-proton collisions at $\sqrt{s} = 8$ TeV.

Clustering algorithm of the PF-reconstruction algorithm: Another clustering algorithm is part of the particle-flow algorithm and individually reconstructs particle showers. Adjacent crystals around a seed which have energies of two standard deviations above electronic noise are grouped to PF clusters. Energy of a crystal can be shared among two or more clusters.

Electron track reconstruction

Tracks of charged particles can be reconstructed using the standard Kalman filter (KF) track reconstruction procedure [86]. As electrons lose a large amount of their energy through bremsstrahlung, the hit-collection efficiency is reduced and track parameters are poorly estimated. Thus, a dedicated tracking procedure is used for electrons. First, in the *seeding* step, SC are matched to track seeds, pairs or triplets of hits in the inner tracker. Matched track-seeds are used as *electron seeds* for electron-track building. Then, tracks are fitted. The performance of seeding greatly affects the reconstruction efficiency. There are two algorithms used for seeding.

ECAL-driven seeding: This algorithm was optimised for prompt electrons in the p_T range relevant for W- or Z-boson decays down to $p_T \simeq 5$ GeV. The SC energy and position is used to estimate the electron trajectory towards the collision vertex. First, electron seeds are selected from SCs, requiring $E_T^{SC} > 4$ GeV, to reduce the number of misidentified seeds. A hadronic veto selection is applied to reduce the background coming from charged hadrons misidentified as electrons: $H/E_{SC} < 0.15$, with H the sum of HCAL tower energies within a cone of $\Delta R = 0.15$ around the electron direction. Then, the electron trajectory is extrapolated towards the inner tracker using the energy-weighted average position of the SC, which is on the helix corresponding to the initial electron-energy, without bremsstrahlung. This back propagation is done for positive and negative charge hypotheses. Intersections of helices with innermost layers or disks are taken as seeding hits. Those seeds are matched to tracker seeds which are formed by combining pairs or triplets of hits with vertices obtained from pixel tracks. For each SC, the SC-predicted hits are compared to hits of tracker seeds within a window in ϕ , and z or in the transverse distance r in the forward regions. Only seeds are taken, whose first two hits are matched to the predictions from SC. Simulated electrons from Z-boson decays have an overall efficiency of ECAL-driven seeding of $\approx 92\%$.

Tracker-driven seeding: This algorithm was optimised for low-momentum electrons ($p_T < 5$ GeV) and electrons within jets and it complements the ECAL-driven seeding for electrons in the barrel-endcap transition region. It was developed as part of the PF algorithm. Tracks reconstructed with the standard KF technique are extrapolated towards the ECAL and matched to an SC. Electron trajectories can be reconstructed accurately with the KF algorithm, when bremsstrahlung is negligible. Hits are collected up to the ECAL. Each KF track has its seeds selected when its direction is compatible with the position of the closest PF cluster that fulfils the matching criterion of $r_{th} < E/p < 3$, with the cutoff r_{th} set to 0.65 for electrons with $2 \text{ GeV} < p_T < 6 \text{ GeV}$ and to 0.75 for electrons with $p_T \geq 6 \text{ GeV}$. When bremsstrahlung is not negligible, the KF algorithm cannot follow the electron trajectory and it either stops collecting hits or collects them with large inefficiencies, leading to a large value of χ^2_{KF} . For this reason, KF tracks with a small number of hits or a large χ^2_{KF} value are refitted using the Gaussian sum filter (GSF) [102]. A multivariate analysis (MVA) is used to select electron seeds from tracker seeds using the number of hits and χ^2_{KF} , the quality of the GSF track χ^2_{GSF} , and the track-SC matching as input.

Electron seeds found by both algorithms are combined, leading to an overall efficiency of $> 95\%$ for simulated electrons from Z-boson decays.

Electron-track building Next, electron seeds are used to build electron tracks. Shorter electron tracks are obtained with the standard KF method than using a dedicated electron building, since bremsstrahlung is neglected and the standard KF algorithm cannot follow the full electron trajectory. The dedicated electron-track building is based on the KF method. For each seed, it proceeds iteratively using track parameters in each layer and information from each successive layer. A Bethe-Heitler function is modelling the energy loss of the electron. Then, it is checked if there is a hit compatible with the predicted one using not too restrictive requirements to follow the trajectory and maintain a good efficiency in case of bremsstrahlung. Trajectory candidates are built if several hits are found which are compatible with the predicted ones. At most one missing hit is allowed per trajectory candidate, but the χ^2 criterion is stricter for such candidates. The number of candidates is limited to 5 for each layer.

Gaussian Sum Filter (GSF) The Gaussian sum filter [102] is used for fitting the electron tracks. Hits found in the electron-track building step are used as input for estimating the track parameters. A model for energy loss of the electron is used which approximates the energy loss in each layer with a mixture of Gaussian distributions, which are weighted according to their associated probability. Track properties can be extracted using the

most probable value for each measurement point. Electron tracks can be followed up to the ECAL and track parameters can be extracted at the surface of the ECAL. With p_{in} the momentum at point of closest approach to the beam spot and p_{out} as the momentum extrapolated to the surface of the ECAL, the fraction of energy lost through bremsstrahlung in the tracker can be defined as:

$$f_{\text{brem}} = (p_{in} - p_{out})/p_{in}. \quad (6.1)$$

This fraction can be used to estimate the electron momentum and also enters into the identification procedure. It is also a perfect tool for estimating the intervening material between collision vertex and ECAL. If the material distribution of detector components is poorly described in simulation, correction factors can be estimated using differences between observed and simulated values for f_{brem} .

Particle-flow clustering: PF electron clusters are build using GSF tracks and PF clusters, independent of the used track-seeding. For each track, PF clusters corresponding to the electron and emitted bremsstrahlung photons are grouped together. The electron PF-cluster can be found by matching tracks at the exit of the tracker to PF clusters. PF clusters corresponding to bremsstrahlung photons are found by extrapolating a straight line to the ECAL tangential to the electron track. Matching PF clusters are added to the electron cluster. Bremsstrahlung photons can convert into e^+e^- pairs. An MVA algorithm is used to find PF clusters corresponding to converted photons, which cannot always be recovered by the other method for bremsstrahlung photons. A dedicated MVA algorithm selects displaced KF tracks and associates them with PF clusters. For ECAL-seeded isolated electrons, PF clusters which match to an SC are added to the PF electron cluster.

Association between track and cluster

Finally, electron candidates are reconstructed by associating GSF tracks to a cluster in the ECAL using not very restrictive track-cluster association criteria. For ECAL-seeded electrons, the SC clusters are used which were reconstructed using either the hybrid or the multi- 5×5 algorithm. A geometrical matching is used as association criterion by comparing the distance in η and in ϕ between the SC and the position of closest approach of the track to the SC. For tracker-seeded electrons, electron PF clusters are used. A global identification variable is defined using an MVA algorithm which combines information on observables of GSF tracks, electron PF clusters, and the association of both.

Converted bremsstrahlung photons can be reconstructed as electron candidates when the photon carries a significant fraction of the initial electron energy and its hits are located close to the expected position of the initial track. When two nearby GSF tracks share the

6. Event reconstruction

same SC, several criteria are applied for selecting the electron track. If both GSF tracks have a different number of missing inner hits, obtained from the intersections between the track trajectory and the active inner layers, the one with the smallest number of missing inner hits is retained. If both tracks have the same number of missing inner hits and just one track has an ECAL-driven seed, this track is chosen, because tracker-driven seeds have a higher probability to be contaminated by track segments from photon conversions. In all other cases, the GSF track with E_{SC}/p closest to unity is chosen, with the track momentum p evaluated at the interaction vertex.

The overall reconstruction efficiency is $\approx 93\%$ for simulated electrons from Z-boson decays.

6.2.2. Electron charge estimation

The electron charge can be estimated using the GSF track curvature. However, bremsstrahlung and photon conversions can affect the measurement of the electron charge. Using two other charge estimates can improve the charge measurement. The first one is based on a KF track which shares at least one hit with the GSF track in the innermost region. The second one is based on the SC position, using the sign of the difference in ϕ between two vectors which originate from the beam spot to the SC position or the first hit of the GSF track.

There are two methods for estimating the electron charge from those three charge estimates. The *majority method* defines the electron charge by the sign which is shared at least between two of these three estimates. The *selective method* requires all three estimates to agree.

6.2.3. Electron momentum estimation

Electron momentum is estimated combining tracker and ECAL measurements. It is sensitive to bremsstrahlung and to conversion photons. SC energy is corrected to account for energy leakages, energy loss and additional energy from pileup. Before clustering, the energy response of individual crystals is calibrated [103] and changes in the response due to radiation are corrected through the ECAL laser-monitoring system [104]. Further corrections in momentum scale and resolution can be estimated using dielectron resonances, which are shown in Figure 6.2. Electrons from Z-boson, J/Ψ and Υ resonance decays are well suited for estimation of correction factors.

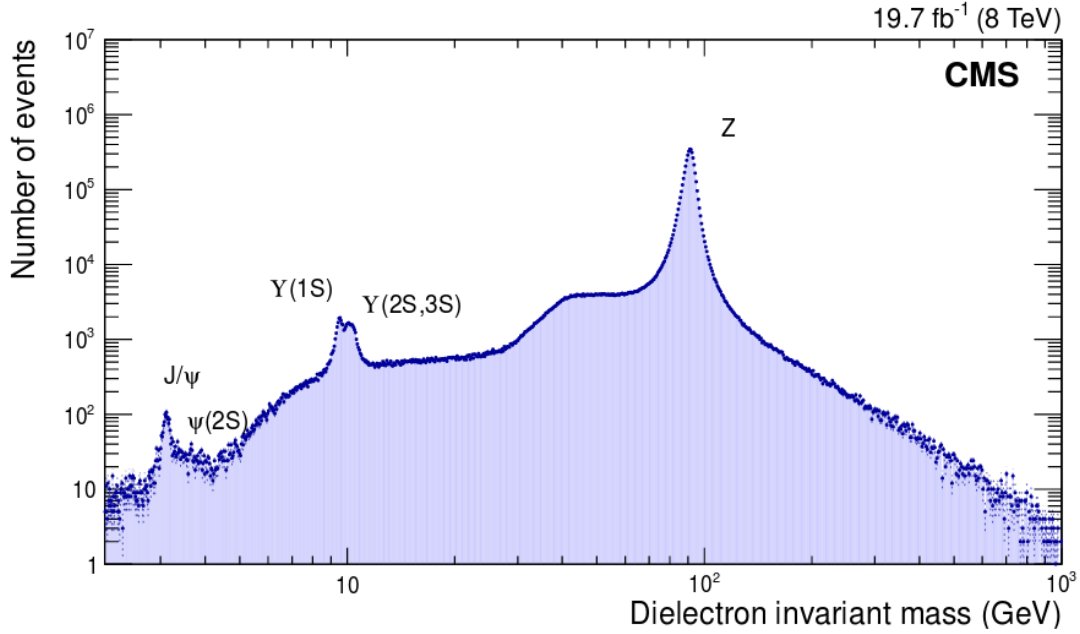


Figure 6.2.: Dielectron invariant mass spectrum for the collision data collected with dielectron triggers in proton-proton collisions at $\sqrt{s} = 8$ TeV [55].

The electron momentum estimate p_{comb} can be improved by using the corrected SC energy and by combining it with the track momentum measurement, which is more precise for energies below 15 GeV and electron near gaps in detectors. At higher energies above typically 250 GeV, using the SC information alone gives a better estimate.

Information of ECAL and tracker is combined by using a regression technique for calculating a weight w which is used for a linear combination of both estimates:

$$p_{\text{comb}} = wp + (1 - w)E_{\text{SC}} \quad (6.2)$$

The following input observables sensitive to bremsstrahlung are used in the regression:

- corrected SC energy and its uncertainty
- track momentum and its uncertainty
- the ratio E_{SC}/p , its uncertainty, and the ratio of the E_{SC} and p uncertainties
- electron class defined by the number of clusters in the SC and the bremsstrahlung fraction in the tracker and its difference to the bremsstrahlung fraction in the ECAL ($f_{\text{brem}}^{\text{ECAL}} = (E_{\text{SC}}^{\text{PF}} - E_{\text{ele}}^{\text{PF}})/E_{\text{SC}}^{\text{PF}}$ with the SC energy and electron-cluster energy)
- position in barrel and endcaps.

6.2.4. Electron identification

For analyses involving electrons, it is important to separate prompt electrons from background sources. These backgrounds can either come from real electrons, such as electrons from photon conversions and electrons in jets from heavy flavour decays, or are fakes coming from jets which are misidentified as electrons. Prompt electrons are isolated, i.e. they have no significant energy flow near their trajectories unlike electrons within jets or fakes.

Different groups of selection observables are used to separate prompt electrons from background. They can be classified as identification (ID) observables, isolation observables and conversion rejection observables. Isolation observables are not included in the electron identification because the requirement on the electron isolation is analysis dependent.

Identification observables

Identification observables can exploit tracking observables, such as separation between electrons and charged hadrons, ECAL information, like the transverse shower shape, or they can compare the ECAL and tracker information, such as geometrical track-cluster matching and SC energy-track momentum matching. Important ID observables are:

- $\Delta\eta_{in}, \Delta\phi_{in}$: spatial matching between track and SC
- H/E_{SC} : hadronic leakage variable
- $\sigma_{in\eta}$: width of the SC in η
- $1/E_{SC} - 1/p_{in}$: with track momentum p_{in} .

Rejection of converted photons

Electrons from converted bremsstrahlung-photons are rejected by making use of the pattern of track hits. First hits of electrons coming from converted photons are in general not located in the innermost layer of the tracker while tracks from prompt electrons start from the beamline, therefore missing hits in the inner layer can be used to reject electrons from photon conversion. Partners of the conversion track can be found by fitting track pairs to a common vertex, as such tracks have the same tangent at the conversion vertex because photons are massless. Finally, requirements on impact parameters IP, defined as distance to the vertex at point of closest approach in the transverse plane, can be applied, with d_{xy} being the transverse IP and d_z being the longitudinal IP. The ratio of the uncertainty of the three-dimensional IP and its value σ_{IP}/IP can be used to reject secondary electrons.

Selections

Analyses have to balance selection efficiency and purity depending on the amount of signal and background. Either a cut-based approach can be used, which uses simple requirements on ID observables, or an MVA technique can be used. A *boosted decision tree* (BDT)¹ can combine the information of those observables into a single discriminator value. Two types of BDTs are defined. For *triggering electrons*, loose identification and isolation requirements have to be fulfilled before using the BDT. Those electrons pass HLT identification. *Not-triggering electrons* do not have to pass such requirements.

Working points (WP) with a specified efficiency for electrons from Z-boson decays are defined. Either a specified selection on the ID observables is given for the cut-based approach, or a selection on the MVA discriminator is given. The *loose* WP has an efficiency of 90%, the *medium* WP 80%, the *tight* WP 70%, and the *very loose* WP used for vetoing electrons has an efficiency of 95%.

6.2.5. Electron isolation

Electron isolation can be used to separate prompt electron from electrons within jets or jets misidentified as electrons. The isolation observable describes the energy flow around the electron trajectory. Electrons in jets tend to have a significant energy flow around their trajectory. There are two different kind of electron isolation observables, which make use of isolation sums, removing contributions from the electron, bremsstrahlung photons and conversion photons from these sums.

Detector-based isolation uses the scalar sum of track momenta reconstructed from the collision vertex, or the sum of energy deposits either in the ECAL or in the HCAL, without taking the momentum or energy deposits from the electron candidate into account. These sums are computed within a cone radius of ΔR of 0.3 or 0.4. This isolation observable is used in the HLT.

PF isolation uses the scalar momentum sums of PF candidates within a chosen cone ΔR around the electron direction. The PF technique combines the information of all subdetectors and avoids double counting. Additional energy depositions originating from pileup is subtracted. The PF isolation is defined as

$$\text{Iso}_{\text{PF}} = \Sigma p_{\text{T}}^{\text{charged}} + \max[0, \Sigma p_{\text{T}}^{\text{neutral had}} + \Sigma p_{\text{T}}^{\gamma} - p_{\text{T}}^{\text{PU}}], \quad (6.3)$$

taking into account charged PF candidates, originating from the vertex, neutral hadrons

¹A BDT is a binary structured tree classifier. The phase space is split into many regions classified as signal or background by repeated yes/no decisions taken on one single variable at a time. The boosting extends this concept to several decision trees. The final decision is taken by a weighted average of the individual decisions [105].

6. Event reconstruction

and photons. Additional energy related to pileup can be estimated using the event-specific average pileup energy density per unit area ρ , which is defined as median of the energy distribution for particles within the area of any jet in the event. An effective area A_{eff} in (η, ϕ) can be defined for each given type of isolation. The effective area is the geometric area of the isolation cone corrected by an η -dependent factor that accounts for the residual dependence of the isolation on the pileup [106]. Energy due to pileup is then defined as

$$p_{\text{T}}^{\text{PU}} = \rho A_{\text{eff}}; \quad (6.4)$$

pileup is discussed in more detail in Section 6.6.

A *relative isolation* can be defined by dividing the Iso_{PF} observable by the electron momentum. Requirements on the isolation can reduce background significantly. This requirement is analysis dependent.

6.3. Muon reconstruction and identification

Muons are the only SM particles which reach the outer parts of the detector and can be identified. A good muon reconstruction and identification is important for many measurements and searches at the Compact Muon Solenoid (CMS) experiment [29]. Since muons leave a clean signature in the detector, searches and measurements often make use of signatures which include muons. The $H \rightarrow ZZ^* \rightarrow 4l$ channel, for example, was among the discovery channels of the Higgs boson, due to its clean signature. Invariant lepton masses and lepton charges are well suited for discriminating Z bosons from other background. Muons can also be used to identify other particles which have a high branching fraction to decay into a muon.

In this section, the reconstruction and identification of muons is discussed [88]. First, the reconstruction of muon is described. Then, algorithms for muon-momentum assignment are introduced. Different muon sources are discussed for a better understanding of the muon identification selections. Identification selections are introduced and algorithms for estimating the muon isolation are described.

6.3.1. Muon reconstruction

First, tracks are reconstructed independently in the inner tracker leading to *tracker tracks* objects and in the muon system leading to *standalone-muon track* objects. In the next step, a matching between both objects is done for muon reconstruction. Two algorithms are used based on these different objects.

Global muon reconstruction (outside-in): Muons are reconstructed using standalone-muon tracks. The matching to tracker tracks is done by comparing parameters propagated onto a common surface for a standalone-muon track and a tracker track. If a match is found, a *global-muon track* is fitted combining hits from both tracks using the Kalman-filter technique [100].

Global muon reconstruction is designed to have high reconstruction efficiency for muons penetrating through a large part of the muon system and typically requires to have segments, i.e. a short track stub made of DT or CSC hits in at least two muon stations. This global-muon fit can improve the momentum resolution at large transverse momenta, $p_T \gtrsim 200$ GeV in comparison to the tracker-only fit.

Tracker muon reconstruction (inside-out): All tracker tracks with $p_T > 2.5$ GeV and with total momentum larger than 2.5 GeV are considered. Tracks are extrapolated to the muon system taking into account the magnetic field, the average energy loss, and multiple Coulomb scattering in the detector material. A track-to segment matching is done using a local (chamber) coordinate system. The local x is the best-measured coordinate in the r - ϕ plane and the local y is the coordinate orthogonal to it. An extrapolated track is matched to a muon segment if the distance between them is less than 3 cm in local x or if the pull for local x is less than 4, with the pull defined as difference between the track position and the segment position, divided by their combined uncertainties.

Tracker muons have a match to at least one muon segment. Since only a single muon segment is required, Tracker muon reconstruction is more efficient than the global muon reconstruction at low $p_T \lesssim 5$ GeV.

About 99% of muons produced in proton-proton collisions are reconstructed either as *Global Muon* or *Tracker Muon*, due to high tracker-track reconstruction efficiency and very high segment reconstruction efficiency in the muon system. Very often they are reconstructed as both and candidates which share same tracker tracks are merged into a single candidate. Muons reconstructed only as standalone-muon tracks have a worse momentum resolution and higher backgrounds coming from cosmic-ray muons than Global and Tracker Muons, therefore they are usually not used in physics analyses.

6.3.2. Muon momentum assignment

Muons are often reconstructed as Tracker Muon and Global Muon, therefore an algorithm is needed to decide which momentum estimate to use. The *sigma switch* is the default algorithm for muon momentum assignment in CMS choosing between the estimates given by the tracker-only fit and by the global fit. In most cases, the tracker-only case is taken, except when both yield muon momenta above 200 GeV and charge-to-momentum ratios q/p that agree within $2\sigma_{p/q}$. In these cases, the muon-momentum estimate of the global fit is chosen.

When the muon is passing through the steel of the magnet return yoke, radiative processes can alter the muon trajectory. This can result in large energy losses producing electromagnetic showers which give rise to additional hits in muon chambers. Muon-momenta estimates at the production vertex can differ from the true value for muons with high momentum, for which muon-momentum estimate of the global fit is used. Specialised algorithms were developed including information from the muon system. The performance of such algorithms can be studied with cosmic-ray muons.

- *Tracker-Plus-First-Muon-Station (TPFMS) fit*: Sensitivity to possible showering deeper in the muon system is reduced by refitting the global muon track using only hits of the innermost muon station containing hits.
- *Picky fit*: Hits of the global-muon track originating from showers are removed before refitting using the hit occupancy of the muon chambers. Hits from chambers which appear to contain hits from showers are only considered for refitting if they are compatible with the extrapolated trajectory based on a χ^2 test.
- *Tune P algorithm*: This algorithm chooses between the tracker-only, TPFMS, and picky fits depending on the goodness of the fit. It starts with the momentum of the picky fit and compares it with the tracker-only fit. Then, the chosen fit is compared with the TPFMS fit.

6.3.3. Muon sources (in simulation)

There are different muon sources. Most muons are produced in proton-proton collisions, but there is also some background coming from cosmic-ray muons and accelerator-induced background (*beam-halo muons*).

There are several categories for muon produced in proton-proton collisions. *Prompt muons* come from W, Z and top-quark decays. Drell-Yan process and promptly decaying quarkonia states also contribute. Other muon sources are *heavy-flavour muons* which arise from bottom hadron, charmed hadron, or tau decays. *Light-flavour muons* arise from decays of light hadrons (π and K) with some originating from decays of particles produced in nuclear interaction of particles with the detector material, or hadrons with mismatched tracker tracks to muon chamber hits. When hadron-shower remnants penetrating through the calorimeter and reach the muon system, these particles are misidentified as muons (*hadron punch-through*). Particles not undergoing nuclear interactions upstream of the muon system (*sail-through*) belong also to this category. A muon can give rise to more than one reconstructed muon candidate. Such candidates can arise from failures of the pattern recognition of the reconstruction, or from patterns which mimic multiple candidates.

Background to collision muons are *cosmic-ray muons*. Such muons can be reconstructed as collision muons or as a pair of oppositely charged muons in the upper and lower detector halves if they pass close to the interaction point. If they are not reconstructed in the tracker, because they are out-of-time or pass too far from the interaction point, they can still be reconstructed as standalone muons. If such a muon is matched to a tracker track, this muon is misreconstructed as a global muon. Cosmic-ray muons can also lead to mismeasured missing transverse energy if they deposit energy in the calorimeters but are not detected in tracker detectors. The ratio of cosmic-ray muons to collision muons increases with increasing muon momentum.

6. Event reconstruction

Another type of background is accelerator-induced. *Beam-halo muons* can overlap with collision events. These muons can be detected by the CSC system, due to its extensive geometrical coverage in the plane transverse to the beam line. They can be identified checking if patterns in the CSC system are consistent with charged particles travelling parallel to the beam line. Beam-halo muons can also induce a large missing transverse energy signal by traversing the barrel or endcap calorimeters at constant ϕ .

For intermediate muon momenta ($30 \text{ GeV} \lesssim p_T \lesssim 100 \text{ GeV}$), prompt muons from W and Z-boson decays dominate. Low-momenta muons ($p_T < 30 \text{ GeV}$) are dominated by semileptonic decays of heavy flavour hadrons, light-flavour decays, and hadron punch-through.

6.3.4. Identification

Physics analyses use muons from different sources and momenta, therefore there are different muon-identification algorithms which are based on various muon-identification variables for background reduction.

Soft Muon selection: This selection is used in quarkonia and B-physics analyses and was optimised for low $p_T < 10 \text{ GeV}$. Tracker Muons are considered for this selection which have at least a matched muon segment in both local coordinates x and y (with $\sigma < 3$). If the segment forms a better match with another tracker track, it is not considered.

Tight Muon selection: Muons from hadron decays are significantly rejected with this selection for a higher purity of prompt muons such as those from W- and Z-boson decays. This selection is used in many physics analyses such in the measurements of inclusive W and Z boson cross sections. It requires Global Muons with a $\chi^2/d.o.f.$ of the global-muon track fit less than 10 and at least one muon chamber hit included to suppress muons from hadron punch-through and hadron decays. The tracker track has to have a match with muon segments in at least two muon stations implying that the muon is also reconstructed as Tracker Muon. This reduces the rate of accidental track-to-segment matches and muons from hadron punch-through. The muon trigger requires segments in at least two muon stations for a meaningful muon-momentum estimation, therefore this requirement makes the selection consistent with events selected by muon triggers. For a good momentum measurement, there have to be more than 10 inner-tracker hits including at least one pixel hit to suppress muons from hadron decays. The transverse impact parameter of the tracker track with respect to the primary vertex $|d_{xy}|$ has to be smaller than 2 mm. This requirement suppresses cosmic muons and further reduces the rate of muons from hadron decays.

Particle-Flow Muon selection: The CMS particle-flow event reconstruction algorithm [91], discussed in Section 6.5, identifies and reconstructs individual particles, such as electrons, muons and hadrons, by combining information from all CMS subdetectors. Particular selection criteria are applied on Global and Tracker Muons depending on the muon environment, for example, muon isolation or energy deposition in the calorimeters. First isolated muons are identified, then muons are identified with a tight and a loose selection. Both selections are optimised for identification of muons in jets with a low rate of misidentifying charged hadrons as muons. Particle-flow reconstruction of leptons was commissioned with leptons from J/Ψ and W-boson decays at $\sqrt{s} = 7$ TeV [107].

All three discussed algorithms are suited for different kind of analyses and each algorithm has its own muon identification efficiency and purity. The $|d_{xy}|$ requirement reduces the cosmic muon background below 10^{-4} for muons with passing tight identification criteria and with $p_T > 10$ GeV. Only a fraction below 5×10^{-5} of collision events is contaminated by beam-halo muons for tight muons with a momentum exceeding 5 GeV. Figure 6.3 shows different muon sources for Tight Muons with a momentum exceeding 20 GeV. At higher muon momentum, prompt muons originating from W- and Z-boson decays dominate, while at low momenta muons from heavy-flavour decays dominate.

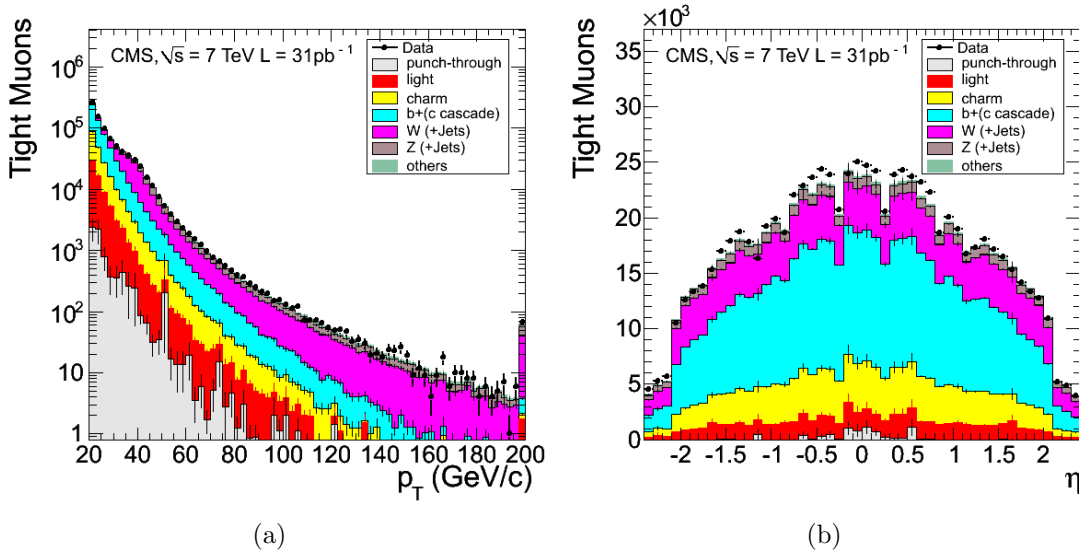


Figure 6.3.: Different muon contributions to Tight Muons in simulation with $p_T > 20$ GeV as a function of (a) transverse momentum and (b) pseudorapidity [88]. Simulated events are compared to observed events recorded in proton-proton collision at $\sqrt{s} = 7$ TeV. Simulated events were generated using PYTHIA 6 [57] with the Z2 tune [108]. Events were recorded in proton-proton collision at $\sqrt{s} = 7$ TeV.

6.3.5. Muon isolation

The energy flow close to a muon, *muon isolation*, can be used to discriminate prompt muons from W- and Z-boson decays from muons produced in heavy-flavour and hadron decays. Different isolation algorithms exist, which sum the energy around a muon in a certain cone of radius ΔR centred on the muon direction. Contributions from the muon are not taken into account.

Tracker relative isolation (I_{trk}^{rel}): Track directions and momenta are computed at the point of closest approach to the nominal centre of the detector. Then, scalar sum of all tracker-tracks momenta reconstructed within $\Delta R < 0.3$ is calculated. The isolation is defined as the ratio of the momenta sum to the muon-track momentum.

Tracker-plus-calorimeters (combined) relative isolation (I_{comb}^{rel}): This algorithm is very similar to the I_{trk}^{rel} algorithm, but it also includes the sum of energies measured in ECAL and HCAL in the numerator sum. Energy deposits within a cone of $\Delta R < 0.3$ are taken into account, not including energy deposits of the muon track.

Particle-flow relative isolation (I_{PF}^{rel}): The sum of charged hadron momenta, transverse energies E_T of all photons and neutral hadrons within a cone of $\Delta R < 0.4$ is calculated using particles reconstructed by the particle-flow algorithm [91]. This sum is then divided by the muon momentum.

Optimisation of cone sizes was done independently for each algorithm. Each algorithm is suited for a different kind of analysis. Measurements of the W and Z boson cross sections use the tracker-plus-calorimeters relative isolation, while searches for heavy resonances use the tracker relative isolation because muons are expected to deposit a significant amount of energy in the calorimeters. Particle-flow relative isolation is used in analyses involving taus, since the particle-flow algorithm is used for global event reconstruction in these analyses.

Methods to include the impact of pileup in the energy sum were developed and are important for higher luminosity runs. One method uses jet areas for estimating the energy in the isolation cone due to particles produced in pileup collisions [109]. The average transverse momentum per unit area ρ added to the event by pileup collisions was measured and a correction factor can be applied on an event-by-event basis. Another method computes a correction factor β using reconstructed tracks and primary vertices [110].

6.4. Lepton efficiency estimation using a tag-and-probe technique

Leptons used in analyses can be selected by using different reconstruction algorithms, identification and isolation criteria. Events can be selected with triggers based on lepton-trigger objects (online), or by matching lepton candidates (offline) with lepton trigger-objects. Each selection step has its own efficiency defined as the probability that the lepton passes the corresponding quality criteria. These efficiencies can be different for leptons reconstructed in collision data and MC simulation, therefore correction factors, scale factors, are estimated and applied on leptons from MC simulation. The efficiency is in general estimated with the *tag-and-probe* technique [89].

This section discusses the estimation of lepton efficiencies [54, 55, 87, 88]. First, the tag-and-probe technique is introduced. Then, application of this technique to electrons and muons is described.

6.4.1. Tag-and-probe technique

Lepton efficiencies can be estimated using the *tag-and-probe* technique [89]. Resonances which decay either into two leptons of same type and have a clean signature can be used for the efficiency measurement. Events are selected by applying strict requirements on one lepton (*tag*) and looser requirements on the other lepton (*probe*). Estimated efficiencies are, to a large extent, insensitive to the requirements on the tag lepton. Criteria applied on the probe lepton have to be less strict than the requirements under study. The efficiency is estimated by the fraction of leptons that pass the selection under study and is determined from fits to the invariant mass distribution of both leptons, which are the decay products of the resonance mentioned before. This distribution includes contribution from signal and background processes as well as from the combinatorial background. Different signal and background models are used to disentangle both contributions. The background component is expected to be described by a falling exponential function with a bias due to the kinematic selection applied on the resonance candidates. Templates from simulation or analytical functions can describe the signal distribution. Several uncertainties are considered. The main systematic uncertainty is related to the used models in the fit. This technique is applied to collision-data events and simulated events. Scale factors can be derived from the efficiency ratio for observed events divided by the efficiency estimated from simulated ones.

6.4.2. Electron efficiencies

Electron efficiencies are estimated by exploiting Z-boson decays to an electron-positron pair with the tag-and-probe technique. Reconstruction, identification, isolation, and trigger efficiencies are estimated in a relative manner, such that the entire efficiency of the selection chain is given as the product of all individual factors. In collision data, events collected with triggers that not bias the efficiency under study are used.

For electrons with $E_T^{\text{SC}} > 5$ GeV, the SC reconstruction efficiency is close to 100%. Triggers requiring one electron and one SC are used for estimating the reconstruction efficiency. The reconstruction efficiency is computed as a function of SC E_T^{SC} and η . Identification efficiencies, isolation efficiencies and trigger efficiencies are computed in bins of electron momentum and SC η .

6.4.3. Muon efficiencies

There are several efficiencies for muons which have to be considered in an analysis. The efficiency to reconstruct a muon in the inner tracker $\varepsilon_{\text{track}}$ was measured to be 99% or higher [111]. Efficiencies for muon reconstruction and identification $\varepsilon_{\text{rec+id}}$, efficiencies for muon isolation ε_{iso} , and trigger efficiencies $\varepsilon_{\text{trig}}$ have to be considered, too. These efficiencies are defined in a relative manner, such that the efficiency of the entire muon selection chain can be calculated as the product of individual factors:

$$\varepsilon_{\mu} = \varepsilon_{\text{track}} \cdot \varepsilon_{\text{rec+id}} \cdot \varepsilon_{\text{iso}} \cdot \varepsilon_{\text{trig}}. \quad (6.5)$$

Efficiencies for prompt muons can be estimated by using a *tag-and-probe* technique to muons from dimuon resonances.

Figure 6.4 shows the invariant mass spectrum of dimuons $M_{\mu\mu}$ collected with a loose double-muon trigger. Several dimuon resonances can be seen including the J/Ψ resonance at $M_{\mu\mu} \sim 3$ GeV as well as the Z-boson resonance at $M_{\mu\mu} \sim 91$ GeV. Branching fractions of both particles to decay into a dimuon pair are small, but muon resonances leave a very clean signature in the detector. The J/Ψ resonance is well suited for muons with low-momenta. For intermediate muon momenta, muons from Z-boson decays can be used.

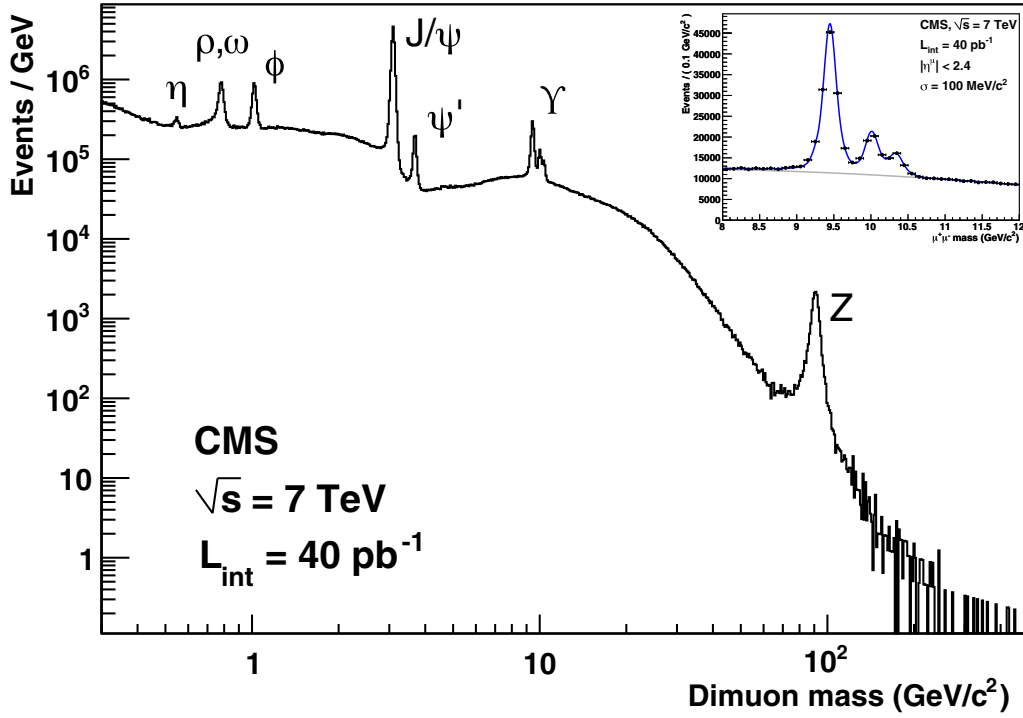


Figure 6.4.: Invariant mass spectrum of dimuons in events collected with a loose double-muon trigger in 2010 [88]. A zoom of the 8 – 12 GeV region is shown on the top right, showing the three $\Upsilon(nS)$ resonances clearly resolved. The mass resolution is about 100 MeV in the entire pseudorapidity range and 70 MeV when both muons are within the range $|\eta| < 1$. Events were collected in proton-proton collisions at $\sqrt{s} = 7$ TeV.

For the Z resonance, a sample of dimuon pairs can be collected by using single-muon triggers. The residual background is subtracted by performing a fit to the invariant dimuon mass spectra $M_{\mu\mu}$ each for passing and failing probe muons. A mass window of $70 \text{ GeV} < M_{\mu\mu} < 110 \text{ GeV}$ is taken for the fit, signal shapes are identical in both fits and appropriate background shapes are used. The efficiency can be computed from the normalisation of the signal shapes in both spectra.

Background and signal modelling, including parameters controlling their shape, the background efficiency and background subtraction contribute to the systematic uncertainty of the efficiency estimated with the tag-and-probe technique.

When estimating reconstruction and identification efficiencies ε_{rec+id} , probes are tracks which were reconstructed using only the inner tracker to have no bias from the muon system. Tight muons, which are defined in Section 6.3.4, with an isolated tag muon are usually used for estimating the isolation efficiency ε_{iso} , since most analyses use tight muons for identification of prompt muons. Trigger efficiencies are estimated using events collected with a single-muon trigger. The identification selection and isolation requirements are applied on the tag and the probe muon. Trigger objects are matched to the reconstructed muons.

Using the tag-and-probe technique with dimuon resonances, an unbiased estimation of muon efficiencies can be done for collision data and MC simulation. Differences in efficiency can be corrected by applying scale factors, defined as the ratio of observed efficiency and efficiency in simulation. Simulated events can be corrected with these scale factors, which are estimated as function of muon momentum and pseudorapidity. Muons with a momentum larger a few GeV have an identification efficiency above 95% and a misidentification probability below 1%.

6.5. Particle-flow event reconstruction

All stable particles in an event can be reconstructed and identified individually with the particle-flow (PF) algorithm [91]. Muons, electrons, photons, charged hadrons, and neutral hadrons can be reconstructed by combining information from all sub-detectors. This leads to an improved measurement of particle momentum, energy, direction and better identification.

Tracks, originating from charged particles, can be reconstructed with large efficiency and small fake rate. Electrons, photons, and to some extent hadrons, deposit energy in the ECAL, whose granularity allows to separate photons from charged-particle energy deposits. Charged and neutral hadrons deposit also energy in the HCAL, which has a worse resolution than the ECAL due to a coarser granularity. Combining ECAL and HCAL information, allows us to detect an energy excess of neutral hadrons on top of the charged-hadron energy and it improves the energy resolution of measured hadrons. Muon tracks are detected in the muon system in the outer part of the CMS detector. Missing transverse energy is determined by the absolute value of the vector sum of transverse momenta of all identified particles and reconstructed jets.

This section discusses the particle-flow algorithm [91]. First, this algorithm reconstructs the fundamental *elements* which are charged-particle tracks and clusters in the calorimeters. Then, these elements are linked to *blocks*. These those blocks are interpreted as particles. The reconstruction and identification performance of jet, tau and missing transverse energy is improved by using the particle-flow algorithm.

6.5.1. Fundamental elements

The PF algorithm uses charged-particle tracks, calorimeter clusters, and muon tracks. Most stable particles have a rather low momentum and since information of all sub-detectors is used by the algorithm, its elements have to be reconstructed with high efficiency and low fake rate even in high particle-multiplicity environments.

Iterative tracking Charged-particle directions at the production vertex can be measured before any deviation due to the magnetic field in the tracker. Since the tracker has a better resolution than the calorimeters, charged-hadron momenta are measured in the tracker. Charged hadrons carry in average two thirds of the energy in jets, therefore tracking is very important for the PF algorithm. The tracking efficiency have to be close to 100% and the tracking fake rate has to be small in order to avoid potential large energy excesses. This is achieved with an iterative-tracking strategy [86]. First, tracks are seeded and reconstructed using tight criteria. This leads to a small fake rate, but a moderate track efficiency. Then, hits which are unambiguously assigned to tracks found in the previous iteration are removed to keep the fake rate low due to reduced combinatorics. In the next iterations, seeding criteria are progressively loosened and the track efficiency increased.

Calorimeter Clustering Electrons, photons and hadrons deposit energy in the calorimeters. Clustering algorithms are used to reconstructed and identify electrons and accompanying bremsstrahlung photons. They detect and measure energy and direction of stable neutral particles and separate their energy deposits from charged hadrons ones. For low-quality tracks or high-momentum tracks of charged hadrons, where an accurately determination of track parameters is not possible, the calorimeter information is used instead for the particle-momentum estimation.

A clustering algorithm has been developed for the particle-flow event reconstruction, which aims to have high efficiencies even for low-energy particles and a good separation of close energy deposits. The clustering is performed separately in each sub-detector, ECAL, HCAL, and ES (pre-shower). It is also performed for the barrel and endcap, and ES first layer and ES second layer separately. No clustering is performed in the HF ($3 < |\eta| < 5$), so each HF cell gives rise to one cluster. First, local cell-energy maxima above a given energy and a given threshold are identified as *cluster seeds*. These *topological clusters* give rise to as many *particle-flow clusters* as seeds. Then, other cells are aggregated to the clusters from this seeds, if they have at least one side in common with a cell already in the cluster. The energy of each cell is shared among all particle-flow clusters according to their cell-cluster distance. Cluster energies and positions are determined iteratively.

6.5.2. Link algorithm

A particle can initiate several particle-flow elements: a charged-particle track, one or several calorimeter clusters, and a muon track. A link algorithm connects those elements for the full reconstruction of each particle and removes double counting. *Blocks* of elements are produced by the algorithm made of directly or indirectly linked elements. The algorithm tries to link each pair of elements in the event. A distance is defined for linked elements to quantify the quality of the link.

6. Event reconstruction

Charged-particle tracks and calorimeter clusters are linked by extrapolating the track from the last measured hit to the two ES layers, to the ECAL at a depth corresponding to an expected maximum of a typical electromagnetic shower, and to the HCAL at a depth corresponding to one interaction length of a typical hadron shower. The track is linked to any cluster if the extrapolated position of the track in the calorimeter is within the cluster boundaries. A link distance is defined as the distance in the (η, ϕ) plane between the extrapolated track position and the cluster position.

Electrons can emit bremsstrahlung photons in the tracker layers. These photons deposit energy tangential to the track at the intersection point between the track and tracker layer. The energy of all bremsstrahlung photons emitted by electrons is collected by extrapolating tangents to tracks at those intersection points to the ECAL. If the extrapolated tangent position is within the cluster boundaries, a link is created.

A link between two calorimeter clusters is established, if the cluster position of the more granular calorimeter is within the cluster envelope in the less granular calorimeter. A link distance is defined as distance between both cluster positions in the (η, ϕ) plane.

Finally, a link between a charged-particle track and a muon track is established if a global fit between both tracks returns an acceptable χ^2 . Such links give rise to *global muons*. The χ^2 defines the link distance. If several global muons can be fitted to a track, only the one with the smallest link distance is kept.

6.5.3. Particle reconstruction and identification

Particles are reconstructed and identified from each block of particle-flow elements.

Particle-flow muons First, each global muon gives rise to a *particle-flow muon* if the combined momentum of the global muon is within three standard deviations of the momentum estimated with the tracker information. Tracks corresponding to particle-flow muons are removed from the block.

Particle-flow electrons Then, *particle-flow electrons* are reconstructed and identified. Electrons can lose energy in the tracker by emitting bremsstrahlung photons leading to short tracks. In a pre-identification step, such electrons are refitted with a Gaussian sum filter [102] described in section 6.2.1. A final identification is performed by combining tracking and calorimeter variables. Tracks and ECAL clusters, including bremsstrahlung photons, of particle-flow electrons are removed from the block.

Remaining elements may give rise to photons, charged and neutral hadrons, or in rare cases to additional muons. Tighter quality criteria are applied to the remaining tracks. The measurement of the track momentum has to be better than the energy resolution expected

for charged hadrons, because otherwise it is measured with more precision independently by the calorimeters.

Neutral particles in the block can be detected by comparing the track momenta with energy deposits in the calorimeters. Expected muon energy deposits in the calorimeter are subtracted. Then, ECAL and HCAL cluster energies have to undergo a calibration procedure, described in section 6.5.4.

If several tracks are linked to one HCAL cluster, the sum of their momenta is compared to the energy deposits in the calorimeter. One track can have several linked HCAL clusters. In this case, only the link to the closest cluster is kept for comparison. ECAL clusters with several tracks can either arise from overlapping photons or from hadronic shower fluctuations. In the later case, the other links should be preserved. Only those links are kept, which leave the total calibrated calorimetric energy smaller than the total charged-particle momentum using an ordered list. In some cases, the calibrated energy of the calorimeter is smaller than the total track momentum. When the difference between both estimates is larger than three standard deviations, a relaxed search for muons and fake tracks is performed leading to possibly additional particle-flow muons. Tracks are removed according to their measured momentum uncertainty as long as the calibrated energy of the calorimeter is larger than the total track momentum.

Particle-flow charged hadrons Each remaining track in the block gives rise to a *particle-flow charged hadron*. Their momenta and energies are taken from their track momentum under the charged pion mass hypothesis. Their momenta are refitted if the calibrated calorimeter-energy is comparable to the track momentum. This combination is relevant for high energies and large pseudorapidities.

Particle-flow photons Photons only deposit energy in the ECAL, therefore ECAL clusters, without a link to a track or a disabled link, give rise to a *particle-flow photon*. Neutral particles can also lead to an excess on top of the deposited energy of charged-particles, which is larger than the energy resolution. Neutral hadrons leave only 3% of the jet energy in the ECAL and 25% of the jet energy is carried by photons, therefore a photon is created with the ECAL energy.

Particle-flow neutral hadrons If the energy of clusters linked to a track is larger than the total charged-particle momentum and the excess is larger than the total ECAL energy, a particle-flow photon is created with the ECAL energy, and a *particle-flow neutral hadron* is created with the remaining part. Remaining HCAL clusters, not linked to a track or with a disabled link, give rise to additional particle-flow neutral hadrons. Neutral-hadron energies are determined by applying the calibration procedure only to the HCAL clusters.

6.5.4. Cluster calibration

The ECAL is already calibrated for photons and electrons and has a different response to hadrons. Hadrons have a non-linear HCAL response and the HCAL is calibrated to 50 GeV pions which do not interact in the ECAL. The ECAL and HCAL cluster energies have to be recalibrated in order to estimate the true hadronic energy deposits in the calorimeters. The calorimeter energy linked to charged hadrons is calibrated according to following function:

$$E_{calib} = a + b(E, \eta)E_{ECAL} + c(E, \eta)E_{HCAL}, \quad (6.6)$$

with the calibrated energy E_{calib} , the ECAL and HCAL energies E_{ECAL} and E_{HCAL} , and the HCAL pseudorapidity η . The true energy E is estimated by either the total charged particle momentum or the total uncalibrated calorimetric energy, choosing the larger estimate. The calibration coefficients a , b , and c are estimated by minimising in each bin of E the sum of differences between the calibrated and true energy of the i^{th} hadron divided by the expect calorimetric energy resolution $\sigma_i^2(E_{calib}^i)$:

$$\chi^2 = \sum_{i=1}^N \frac{(E_{calib}^i - E^i)^2}{\sigma_i^2(E_{calib}^i)}. \quad (6.7)$$

The sum extends over all events, separately for the barrel and end-cap and distinguishing if the hadrons only leave their energy in the HCAL or ECAL, or in both calorimeters. The energy resolution is determined iteratively in each bin of E as Gaussian sigma of the $(E_{calib}^i - E^i)$ distribution. The constant correction a can be understood as a constant correction for the thresholds applied to the calorimeter cell-energies in the clustering algorithm.

6.6. Jet reconstruction

Jets are the experimental signature of energetic quarks and gluons produced in high-energy processes. Due to QCD confinement which only allows for colourless states, quarks and gluons are not directly observable. They fragment to a collimated bunch of hadrons flying roughly in the same direction, which can be simplified as a jet.

First, different algorithms for jet reconstruction are described. The anti- k_t algorithm [112], which assigns particles to a jet, and other similar clustering algorithms are introduced. Finally, jet energy scale and resolution are introduced and different consecutive correction factors are discussed. More information about jets in CMS can be found elsewhere [92].

6.6.1. Jet reconstruction

At hadron-hadron colliders, quarks and gluons are dominantly produced. Due to the high luminosity at the LHC, fast algorithms are needed for the reconstruction of jets. These algorithms have to be infrared and collinear (IRC) safe, i.e. soft emissions and collinear splittings must not alter the hard jets. This ensures cancellation of divergence in higher order calculations. Furthermore, background contributions coming from underlying events and pileup are also suppressed.

Jet algorithms need reconstructed particles as input. Charged hadrons can be detected in the tracking detector and calorimeters, while neutral hadrons can be only detected in the calorimeters. One possibility is to use only energy deposits in the calorimeters for jet reconstruction. As charged hadrons lead to tracks and hadrons can decay further to leptons and photons, reconstruction of all involved particles can help to improve the jet reconstruction. Clustering algorithms classify calorimeter energy-depositions or reconstructed particles into jets.

Calorimeter (CALO) jets are reconstructed from energy deposits in the calorimeters alone. Information of one or more HCAL cells and geometrical corresponding ECAL crystals are combined to a calorimeter tower. A four-momentum is associated with each tower deposit above a certain threshold taking the direction as seen from the interaction point.

Particle-flow (PF) jets are reconstructed by clustering four-momentum vectors of PF candidates, which have been reconstructed using the PF technique described in Section 6.5. Using PF candidates for jet reconstruction greatly improves the jet momentum and spatial resolution with respect to calorimeter jets [92]. Tracking information and the high granularity of the ECAL improve the energy resolution through an independent measurement of charged hadrons and photons inside a jet. Both together are responsible for about 85% of the average jet energy. As PF jets have a better resolution than calorimeter jets, they are mostly used in CMS analyses.

Jet-clustering algorithm classify particles into jets. There are different types of clustering algorithms. *Cone algorithms* use a top-down approach by finding coarse regions of energy flow in conical regions in $(\eta-\phi)$ space, while *sequential recombination algorithms* use a bottom-up approach by combining particles starting from closest ones using a certain distance measure. One of such sequential recombination algorithms is the anti- k_t algorithm [112], which is commonly used for jet clustering in CMS.

Jets are built by applying the clustering procedure to either calorimeter towers, or PF candidates. Finally, jets used in analysis are required to satisfy basic identification criteria, *Jet ID*. *Simulated particle level jets* can be built by applying the clustering procedure to all stable (lifetime $c\tau > 1$ cm) particles excluding neutrinos. This defines the response of the detector to jets in a way that is experimentally accessible and significantly reduces response differences between heavy-flavour (c,b) and light-quark (u,d,s) or gluon jets caused by neutrinos in semileptonic decays of heavy-flavour hadrons.

Jets are commonly reconstructed in CMS using PF candidates and the anti- k_t algorithm. A jet distance parameter of $R = 0.5$ is used for $\sqrt{s} = 8$ TeV proton-proton collisions, while a smaller distance parameter of $R = 0.4$ is used for $\sqrt{s} = 13$ TeV proton-proton collisions. Smaller jet radii reduce the amount of underlying event and pileup contributions, which have increased for $\sqrt{s} = 13$ TeV with respect to $\sqrt{s} = 8$ TeV.

6.6.2. The anti- k_t algorithm

In the past, inclusive jet finding algorithms, such as the k_t and Cambridge/Aachen algorithm, have been used. Both are sequential recombination jet algorithms parameterised by the power of the energy scale in their distance measure. Unlike standard cone algorithms, both algorithms are IRC safe. But they suffer from their sensitivity to non-perturbative effects, such as hadronisation and underlying event contamination, which influence the boundaries of reconstructed jets. The anti- k_t algorithm [112] is very similar to both algorithms, but behaves like an idealised cone algorithm for jets with only soft radiation.

The k_t , Cambridge/Aachen, and anti- k_t algorithm introduce distance measures d_{ij} between two particles or pseudojets i and j , and d_{iB} between i and the beam (B). The clustering starts by identifying the smallest of the distances. If the smallest distances is the distance between two entities i and j , both are combined, otherwise i is called a jet and is removed from the list of entities. Distances are recalculated and this procedure is repeated until no entities are left.

The distance measures can be generalised to [112]:

$$d_{ij} = \min(k_{ti}^{2p}, k_{tj}^{2p}) \frac{\Delta_{ij}^2}{R^2}, \quad (6.8a)$$

$$d_{iB} = k_{ti}^{2p}, \quad (6.8b)$$

where $\Delta_{ij}^2 = (y_i - y_j)^2 + (\phi_i - \phi_j)^2$, and k_{ti} , y_i and ϕ_i are respectively the transverse momentum, rapidity and azimuth of particle i . There are two parameters, the radius parameter R , and the parameter p which governs the relative power of the energy versus geometrical scales (Δ_{ij}). The k_t algorithm has a parameter $p = 1$ and the inclusive Cambridge/Aachen algorithm corresponds to $p = 0$. The anti- k_t algorithm has a parameter of $p = -1$.

Due to the negative power in case of the anti- k_t algorithm, softer particles tend to cluster with hard particles before they cluster among themselves. Hard particles with no other hard particle within a distance $2R$ accumulate all soft particles within a circle of radius R leading to a perfectly conical jet. If there is another hard particle within $R < \Delta_{ij} < 2R$, two hard jets are reconstructed, but it is not possible for both to be perfectly conical. Depending on the transverse momenta of both jets, the harder jet is conical and the other one misses the part overlapping with the other jet, or neither jet is conical. Soft particles do not modify the shape of the jet, while hard particles do.

The jet size and the jet area [113] determine the susceptibility of a jet to soft radiation. Larger jet radii allow to capture enough of the hadrons and their decay products to calculate the jet energy. However, a smaller jet radius reduces the amount of underlying event and pileup contributions. Two types of jet areas can be defined. The passive area a measures the jet susceptibility to point-like radiation sources, while the active area A measures the jet susceptibility to diffuse radiation, such as underlying event activity and pileup. Since boundaries of anti- k_t jets are unaffected by soft radiation, the passive area is given by the area of a circle with radius equal to the radius parameter R . As soft particles only cluster among themselves after clustering with hard ones, passive and active jet area are identical:

$$A_{anti-k_t, R} = a_{anti-k_t, R} = \pi R^2. \quad (6.9)$$

6. Event reconstruction

The jet area of anti- k_t jets is independent of the jet momentum. Configurations with several neighbouring hard particles lead to a jet area which deviates from πR^2 .

Adding soft events (underlying events and pileup) contributions to an event can change the subset of reconstructed particles. The added soft particles have a density ρ of transverse momentum per unit area. The probability of soft particles changing the subsets, is suppressed by the anti- k_t algorithm. The additional contributions to the jet energy due to soft events can be subtracted from the reconstructed jets..

6.6.3. Jet energy scale and resolution

As all other experimental-reconstructed objects, jets have to be calibrated for a correct jet energy scale (JES), which would have been obtained, on average, if all particles inside the jet cone would have been measured perfectly. The jet energy resolution (JER) is calibrated as well using jet energy corrections (JEC). A detailed description on how JEC are obtained can be found elsewhere [92].

Figure 6.5 gives an overview on the consecutive stages of the JEC for data and MC simulation. The JEC are calculated using a detailed MC simulation of the CMS detector, and are afterwards adjusted for data. First, it is corrected for the offset energy coming from pileup, multiple proton-proton collisions in the same and adjacent beam crossings. Then, the detector response to hadrons is calibrated. Finally, residual differences between data and MC simulation are corrected as a function of the jet pseudorapidity η and the transverse momentum p_T . The jet momentum is corrected to so-called particle-level jets clustered from stable and visible final-state particles. The JEC are extracted for jets with $p_T > 10$ GeV and $|\eta| < 5.2$. Jet-flavour corrections are optional and derived from MC simulation. The simulated JER default corrections are provided for the QCD dijet flavour mixture as a function of p_T and η . They are estimated after applying the JEC.

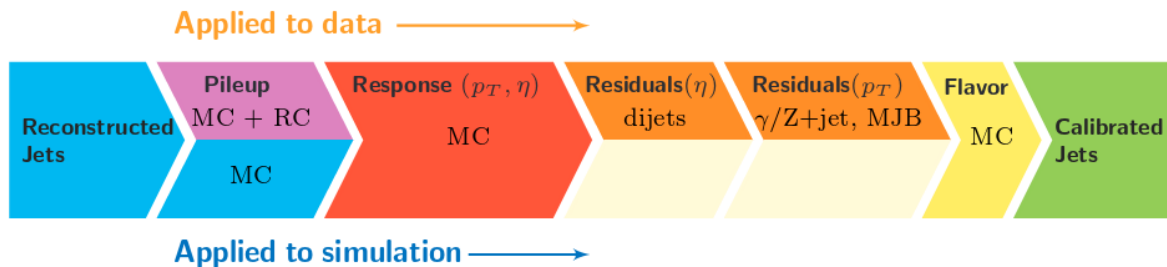


Figure 6.5.: Consecutive stages of JEC, for data and MC simulation. All corrections marked with MC are derived from simulation studies, RC stands for random cone, and MJB refers to the analysis of multijet events [92].

Pileup offset corrections

Multiple proton-proton collisions in the same and adjacent beam crossings (pileup) lead to an offset in the jet energy. In-time pileup (IT PU) occurs when multiple proton-proton collisions occur within the same bunch-crossing as the primary hard interaction. Additional tracks are produced and additional energy is deposited in the calorimeters. Collisions occurring in the previous and subsequent beam crossing can also contribute to additional calorimetric energy due to the finite signal decay time in the calorimeters. This contribution is called out-of-time pileup (OOT PU). Both pileup contributions lead to additional contributions to the jet energy and jet momentum referred to as *pileup offset*. Subtracting pileup contributions improves the detector resolution and leads to a more accurate JES [92].

The amount of pileup in an event can be estimated by the number of good quality reconstructed primary vertices N_{PV} . Interaction vertices are reconstructed using track information only.

The diffuse offset energy density ρ is another estimator for the amount of pileup in an event. It is defined as the median of jet momenta $p_{T,i}$ divided by their area A_i :

$$\rho = \text{median}(p_{T,i}/A_i). \quad (6.10)$$

Another pileup estimator is the average number of pileup interactions per crossing μ , which is obtained by multiplying the instantaneous luminosity, measured using luminosity monitors, with the effective minimum bias cross section σ_{MB} for the corresponding centre-of-mass energy. Simulated events are reweighted to match the distribution of the average number of pileup (μ) in data.

Figure 6.6 shows the number of reconstructed primary vertices N_{PV} and the offset energy density ρ after *pileup reweighting*. It can be seen that N_{PV} is strongly correlated with μ , while ρ is sensitive to the amount of energy deposited by each interaction, and therefore less correlated to μ . Thus, pileup reweighting does not remove all differences between data and MC in the ρ distribution.

6. Event reconstruction

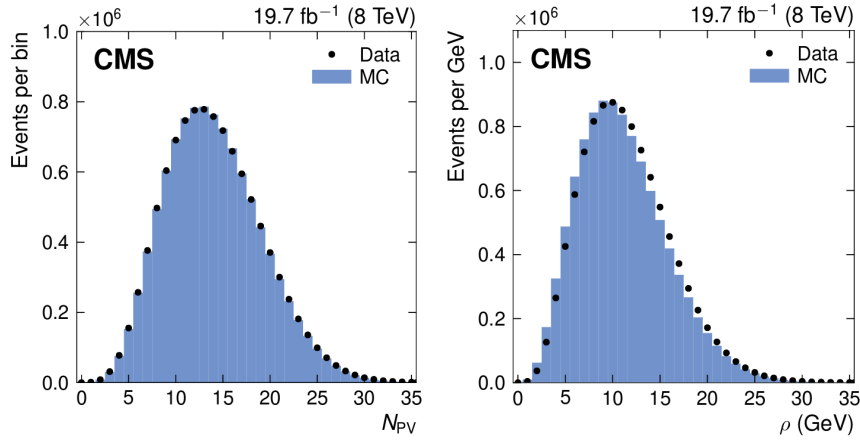


Figure 6.6.: Comparison of data (circles) and PYTHIA 6.4 simulation (histograms) for the distributions of the number of reconstructed primary vertices N_{PV} (left), and of the offset energy density ρ (right) [92].

Additional calorimetric energy depositions and extra tracks originating from IT PU can be reduced by identifying pileup vertices and removing their charged PF candidates before jet clustering. They have to be unambiguously associated with pileup vertices. This method is called *charged-hadron subtraction* (CHS). The primary vertex of the hard interaction is the leading primary vertex defined as vertex with the highest sum of squared transverse momenta of tracks ($\sum |p_T^{track}|^2$) associated with it. Primary vertices are required to pass quality requirements. They have to be compatible with the luminous region and their minimum number of degrees of freedom

$$N_{dof} = -3 + 2 \sum_{i=1}^{nTracks} w_i, \quad \text{and } w_i \in [0, 1] \quad (6.11)$$

of their adaptive vertex fit with corresponding weight w_i has to be $N_{dof} > 4$, which corresponds to at least four tracks. A track is associated only with one vertex if $\chi^2/N_{dof} < 20$ for this vertex.

After applying pileup reweighting, the offset energy can be subtracted from the jets. This can be done using the jet area method which multiplies the effective area of jets by the average energy density in the event. In CMS, a *hybrid jet area method* is used, which combines the η -independent jet area method with an η -dependent average offset correction versus N_{PV} . The pileup offset corrections C_{hybrid} dependent on the uncorrected jet p_T , jet pseudorapidity η , jet area A_j , and the p_T offset density ρ . Parameters of the hybrid jet area method in data are determined using the random cone (RC) method in events without contributions from the hard scattering were the main contributions to jet energies come from noise and pileup, *zero-bias events*. It is assumed that the noise contribution is negligible in comparison to the contribution originating from pileup. The

RC method reconstructs jets in the events by clustering particles in randomly placed cones, mapping all the (η, ϕ) space. The average energy offset due to pileup can be estimated using the average p_T of jets measured from the RC method. In simulation, the pileup offset is calculated by reconstructing the same events with and without pileup overlay and match the jets between both samples to the same particle jet within a distance less than $\Delta R < R/2$. An offset data/simulation scale factor is estimated from zero-bias data and simulation using the RC method. Different working points for the average energy offset density $\langle \rho \rangle$ are used due to the slight differences between data and simulation, as shown in Figure 6.6.

Response corrections

The detector response to hadrons is calibrated using response corrections. These corrections are derived from a simulated QCD multijet sample after applying pileup reweighting and pileup offset corrections. Simulated events are used to cover also corners of phase space, jets with very low or very high transverse momentum, or events with very low or high amount of pileup. The events are generated with a flat p_T spectrum and reweighted as function of $\hat{p}_T^{-4.5}$ with $\hat{p}_T$ being the transverse momentum of the generated $2 \rightarrow 2$ hard process. This reweighting recovers the original p_T spectrum. Fragmentation of initial quarks and gluons was simulated for $\sqrt{s} = 8$ TeV using PYTHIA 6.4 and the Z2* tune [114, 115]. The CMS detector was simulated using GEANT4.

Particle-level jets are matched to the closest reconstructed jet within half of the jet distance parameter R . The simulated particle response R_{ptcl} is defined as ratio of arithmetic means of matched reconstructed and particle-level jets transverse momenta binned in particle-level p_T ($p_{T,\text{ptcl}}$) and reconstructed η :

$$R_{\text{ptcl}}(\langle p_T \rangle, \eta) = \frac{\langle p_T \rangle}{\langle p_{T,\text{ptcl}} \rangle} [p_{T,\text{ptcl}}, \eta]. \quad (6.12)$$

The jet response is sensitive to the underlying detector calibration of the calorimeters. Including tracking information reduces this effect, but since tracking performs worse for higher momenta, the sensitivity to the detector calibration increases with p_T for PF jets.

Residual corrections for data

Residual data/simulation scale factors for JEC account for any residual differences between data and simulation after applying pileup reweighting, pileup offset energy corrections and correcting jets to simulated particle response. Residual corrections are calculated using dijet, $Z(\rightarrow \mu\mu) + \text{jet}$, $Z(\rightarrow ee) + \text{jet}$, $\gamma + \text{jet}$, and multijet events. The p_T balance between a jet, to be calibrated, and a reference object is exploited. Differences in the jet

6. Event reconstruction

energy scale between particle-level jets and reconstructed jets lead to an imbalance at reconstructed level. The *p_T-balance method* evaluates the jet response by comparing the reconstructed jet directly to the transverse momentum of the reference object, while the missing transverse momentum projection fraction method, *MPF method*, considers the response of the whole hadronic activity in the event, recoiling versus the reference object.

Relative η -dependent corrections calibrate the forward regions relative to $|\eta| < 1.3$ using dijet events. The "tag" jet is required to have a pseudorapidity of $|\eta| < 1.3$ and the "probe" jet is unconstrained in η . Corrections are derived using the MPF method, since this method is less sensitive to ISR and FSR biases than the p_T-balance method. Then, absolute JES corrections for the jet p_T are derived using Z + jet and γ + jet events.

Jet flavour corrections

The jet flavour corrections are optional. They are not used for the analyses described in this thesis. These corrections are used to correct for dependencies in the jet response for different jet flavours. These differences arise mainly from variations in the jet fragmentation energy and in the particle compositions. A softer fragmentation results in more particles outside the detector acceptance. The response is mostly affected by differences in the neutral hadron fraction. Gluons have the lowest response, because they fragment into the largest number of soft particles. Light-quark jets, originating from up or down quarks, have the highest response. Heavy-quark jets, originating from c quarks or b quarks, have a smaller response than light-quark jets, since they have additional soft particles from heavy-flavour hadron decays.

A jet flavour can be assigned to a jet by matching a particle-level jet to the nearest generator-level parton within half of the jet distance parameter R. Only partons which belong to the hard scattering matrix are considered, therefore the *parton flavour* gives the jet flavour. Jets without a parton belonging to the hard scattering matrix have an undefined flavour. The response is parameterised as a function of p_T and η . JEC flavour-uncertainties can be estimated comparing PYTHIA and HERWIG predictions.

Jet-p_T resolution

The jet-p_T resolution is relatively poor compared to electrons or muons. Particle-level JER is defined as the width, estimated with a Gaussian fit, of the distribution $p_{T,\text{reco}}/p_{T,\text{ptcl}}$, with transverse momentum $p_{T,\text{reco}}$ of the reconstructed jet, and the transverse momentum $p_{T,\text{ptcl}}$ of the particle-level jet. JEC are applied before deriving the JER corrections. Dijet and γ + jets events are used for determining the jet-p_T resolution. In data, a dijet asymmetry method is used to measure the jet-p_T resolution.

6.7. Missing transverse energy

Neutrinos traverse the detector without producing a direct response in the detector elements, as the probability to interact with the detector material is negligible. As the total momentum is conserved, the presence of such particles can be detected by the imbalance of the total momentum. Protons are composite particles, therefore it is not possible to predict the total momentum component along the beam line. The total momentum component in the plane perpendicular to the beam direction is zero, thus transverse momenta $\vec{p}_T$ and transverse energies E_T are used for estimating the *missing transverse momentum*, $\vec{E}_T$. Its magnitude is called *missing transverse energy*, E_T [93]. Missing transverse energy is an important observable for discriminating events with neutrinos, such as leptonic W-boson decays, from events without E_T , like multijet events. New weakly interacting, long-lived particles would lead to large missing transverse energy, therefore E_T is an important observable for searches for such particles.

In the following, the reconstruction of E_T is described. Corrections due to detector effects are discussed. Finally, it is described how scale and resolution of the reconstructed missing transverse energy are estimated.

6.7.1. Reconstruction of missing transverse energy

The missing transverse energy in an event is given by the negative sum of all transverse momenta of all final-state particles reconstructed in the detector. There are different algorithms available for reconstruction of E_T .

Calo E_T is calculated using energies contained in calorimeter towers and their direction to define pseudo-particles. Energy deposits below noise thresholds are excluded. Muons deposit only a few GeV on average in the calorimeters, which leads to negligible energy deposits in the calorimeter. Their transverse momentum is included in the calculation of Calo $\sum E_T$, which is the associated scalar sum of transverse energies of the calorimeter towers and muon transverse momenta.

TC E_T corrects Calo E_T by including tracks reconstructed in the inner tracker. Track momenta, which have been reconstructed in the inner tracker, are included in the E_T calculation.

6. Event reconstruction

The expected calorimetric energy deposits of each track are removed. Predicted energy deposits for charged pions are used for all tracks, which have not been identified as electron or muon. These predicted energy deposits are estimated from simulation of single pions in intervals of p_T and η using an extrapolation of the CMS magnetic field for the determination of the expected position of the track, when traversing the calorimeter. For high momentum tracks with $p_T > 100$ GeV, no correction is applied, as their energy is better measured by the calorimeters. The measured momentum is taken into account for low momentum tracks, $p_T < 2$ GeV, assuming no response from the calorimeters.

PF $\cancel{E}_T$ is calculated using a particle-flow technique described in Section 6.5. PF $\cancel{E}_T$ is the scalar of the associated vector sum calculated from the transverse energies of all reconstructed PF candidates in the event.

6.7.2. Corrections for reconstructed missing transverse energy

Reconstruction of $\vec{\cancel{E}}_T$ is sensitive to mismeasurements of the particle momenta and particle misidentification. It has to be distinguished between genuine $\cancel{E}_T$, produced by weakly interacting particles, and artificial $\cancel{E}_T$ from detector resolution effects. Underestimation of $\cancel{E}_T$ is mainly due to the displacements of particles with a small momentum due to the magnetic field and the non-linearity of the calorimeter response for neutral and charged hadrons. Calo $\cancel{E}_T$ is affected most by these effects. Corrections can remove this bias in the $\cancel{E}_T$ scale. These corrections rely on the fact that $\vec{\cancel{E}}_T$ can be factorised into contributions from jets, isolated high p_T electrons, muons, and *unclustered energies*, which is defined as difference between $\vec{\cancel{E}}_T$ and the negative of the vector sum of the transverse momenta of the other objects. Isolated photons, electrons and muons are assumed to require no scale correction, whereas jets can be corrected to particle level using jet energy corrections.

Type-I corrections propagate the jet energy corrections to $\vec{\cancel{E}}_T$ for all jets that have less than 0.9 of their energy in the ECAL. Their corrected transverse momentum $\vec{p}_{T,\text{corr}}$ is required to be larger than 20 GeV for Calo $\cancel{E}_T$, and has to exceed 10 GeV for PF $\cancel{E}_T$. The jet energy corrections are propagated to $\vec{\cancel{E}}_T$ by using following equation:

$$\vec{\cancel{E}}_{T,\text{typeI}} = \vec{\cancel{E}}_{T,\text{uncorr}} + \sum_i \vec{p}_{T,\text{uncorr}}^i - \sum_i \vec{p}_{T,\text{corr}}^i - \sum_i \vec{\mathcal{O}}_{\text{RC}}^i, \quad (6.13)$$

where $\vec{\cancel{E}}_{T,\text{typeI}}$ and $\vec{\cancel{E}}_{T,\text{uncorr}}$ are type-I corrected and uncorrected $\vec{\cancel{E}}_T$, respectively, $\vec{p}_{T,\text{uncorr}}$ is the uncorrected jet p_T , and using the average offset due to pileup $\vec{\mathcal{O}}_{\text{RC}}$, as obtained with the RC method [92].

Type-II corrections correct the remaining soft jets below these thresholds, and *unclustered energy*, energy deposits which have not been not clustered in any jet.

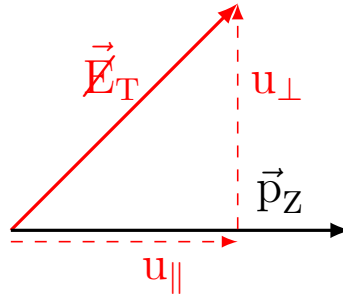
Both corrections have to be applied for Calo $\cancel{E}_T$ while usually only type-I corrections are applied for PF $\cancel{E}_T$.

$\cancel{E}_T$ filters Detector malfunctions, particles traversing poorly instrumented regions of the CMS detector, cosmic-ray particles and beam-halo particles may result in artificial $\cancel{E}_T$. Small anomalous signals could be removed by cleaning the events of such signals, while *filters* tag such an event as not suitable for use in physics analysis for large contributions.

6.7.3. Scale and resolution of missing transverse energy

Isolated photons or Z bosons which decay into $\mu^+\mu^-$ or e^+e^- can be produced in hard parton-parton collisions such as: $qg \rightarrow q\gamma$, $q\bar{q} \rightarrow Z$, $qg \rightarrow qZ$, or $q\bar{q} \rightarrow gZ$. In such events, there is no genuine $\cancel{E}_T$. The detector response can be probed, and the scale and resolution of $\cancel{E}_T$ can be measured by comparing momenta of the Z boson to the induced $\cancel{E}_T$, when removing the vector boson. Measurements of the hadronic products of the hard collisions are influenced by effects due to jet energy scale corrections, underlying event activity, pileup, detector noise, and finite detector acceptance.

The sum of the transverse momentum of the vector boson, q_T , and of the hadronic recoil, u_T , is zero. When the vector boson is removed, the sum of all transverse momenta has to be zero. From this it follows that the induced $\cancel{E}_T$ is the negative of the hadronic recoil, which can be split into a component parallel $u_{\parallel}$ and perpendicular $u_{\perp}$ to the direction of the vector boson with momentum $\vec{p}_Z$.



The response is given by the mean value of

$$\langle u_{\parallel} \rangle / q_T, \quad (6.14)$$

and should be ideally equal to unity as function of q_T . The resolution can be estimated by using the Root-mean-square (RMS) of the difference between $u_{\parallel}$ or $u_{\perp}$ and their mean value, respectively, after correcting for their response.

Using PF candidates for $\cancel{E}_T$ reconstruction improves the $\cancel{E}_T$ resolution with respect to using tracks and calorimetric energy deposits alone.

6.8. Identification of b-quark jets

Jets originating from bottom quarks are present in many physics processes such as top-quark decays, where the top quark decays almost solely into a W boson and a b quark. Accurately identifying jets arising from b-quark hadronisation helps to reduce the overwhelming background from processes involving jets from gluons (g) and light-flavoured quarks (u,d,s), and to a lesser amount jets originating from c-quarks. Furthermore, the identification of b-quark jets helps to suppress the background coming from top decays in searches for decays which do not involve b-quark jets.

During hadronisation, bottom hadrons are formed out of b quarks. The properties of these hadrons can be used to identify b-quark jets, since they have relatively large masses, long lifetimes, and carry a large energy fraction of the bottom quark. Their decay particles also have a hard momentum spectrum. Charmed hadrons have similar properties and, therefore are the main background for the identification of b-quark jets.

The CMS detector with its precise charged-particle tracking is well adapted for the task of b-quark jet identification. Several algorithms exist for identifying b-quark jets [94]. They make use of impact parameters of charged-particle tracks and reconstructed vertices. A commonly used algorithm in CMS is the Combined Secondary Vertex (CSV) algorithm, which uses the lifetime and mass properties of the b hadrons to discriminate between jets originating from b quarks and light-parton (u,d,s,g) jets. Each algorithm yields a single discriminator value for each jet.

First, the reconstructed objects used in b-quark jet identification are defined. Then, the CSV algorithm is described, followed by a description how efficiencies of this algorithm can be measured. These parts are based on b-quark jet identification for LHC Run I [94, 95]. There are some changes in b-quark jet identification for LHC Run II [96], which are listed in the last part of this section.

6.8.1. Reconstructed objects used in b-quark jet identification

After the hadronisation of b quarks, the formed bottom-hadrons decay and the tracks of their charged decay-products can be reconstructed. Due to the large life time of the bottom hadrons, the primary vertex, originating from the hard interaction, and the secondary vertex, originating from the hadron decay, can be reconstructed separately. Algorithms for b-quark jet identification can use track impact parameters, defined as the distance between the point of closest approach of the track to the primary vertex, and the secondary vertices to distinguish b-quark jets from light-parton jets.

Reconstructed objects Jets are clustered from particle-flow objects using the anti- k_T clustering algorithm with a distance parameter of 0.5. Particle-flow jets are used as default, but also jets clustered from other reconstructed objects can be used for the b-quark jet identification algorithms. The jets are required to have a transverse momentum exceeding 20 GeV and a pseudorapidity $|\eta| < 2.4$.

Charged-particle trajectories are reconstructed in the CMS tracking system [86]. For b-quark identification, a sample of well-reconstructed tracks of high purity is required to reduce the amount of fake tracks, coming from a combination of unrelated hits or badly reconstructed tracks.

These reconstructed tracks are clustered based on the z coordinate of their closest approach to the beam line to select primary vertex candidates. The vertex position is estimated using an adaptive vertex fit [101] with a sample of tracks compatible with originating from the interaction region. The candidate with the highest $\Sigma(p_T^{\text{track}})^2$ with the track transverse momentum p_T^{track} is selected as primary vertex.

Track impact parameter based observables Additional requirements are imposed on the track selection for track impact parameter based observables. Tracks are required to have a transverse momentum of $p_T^{\text{track}} > 1$ GeV. At least 8 hits in the tracking system have to be associated with the track and at least 2 hits in the pixel system. For the global track fit, $\chi^2/\text{n.d.o.f} < 5$ is required, where n.d.o.f stands for the number of degrees of freedom in the fit. The impact parameter of a track is defined as distance between the track and the primary vertex at the point of closest approach in the transverse plane. The longitudinal impact parameter d_z , defined as distance along the beam axis, has to be smaller than 17 cm and the transverse impact parameter d_{xy} , the distance in the transverse plane has to be smaller than 0.2 cm to reduce the background coming from long-lived particles, such as neutral kaons. Tracks are associated to a jet in a cone $\Delta R < 0.5$ around the jet axis, which is defined by the primary vertex and the direction of the jet momentum. The distance of a track to the jet axis is defined as the distance between the point of closest approach to the axis. It is required that the distance of a track to the jet axis is less than

700 μm to reject tracks from pileup. Additionally the point of closest approach has to be within 5 cm of the primary vertex.

The impact parameter of a track with respect to the primary vertex, IP, can be used to distinguish decay products of *b*-hadrons from prompt tracks. It is calculated in three dimensions due to the excellent resolution of the pixel detector. The IP resolution depends on track momentum and pseudorapidity. The impact parameter has the same sign as the scalar product of the vector pointing from the primary vertex to the point of closest approach with the jet direction [94]. Decay particles travelling along the jet axis tend to have positive IP values, while prompt tracks can have positive and negative IP values. The impact parameter significance S_{IP} , defined as the ratio of the IP to its estimated uncertainty, can be used as observable for *b*-quark jet identification.

Secondary vertices based observables For the reconstruction of the secondary vertices coming from bottom and charmed hadron decays, a different track selection is applied. Tracks have to be within a smaller jet cone of $\Delta R < 0.3$. The distance of closest approach of a track to the jet axis has to be less than < 0.2 cm. An adaptive vertex fit is used for finding vertex positions. The vertex finding starts with the tracks passing this track selection and reconstructs the vertices iteratively. In the first iteration, primary vertices are reconstructed. Then, the secondary vertices originating from particle decays are reconstructed.

The presence of a secondary vertex and its kinematic variables can be used to discriminate jets originating from *b*-quark jets from other jets. Such variables are the flight distance and the flight direction, using the vector between primary and secondary vertex. The flight distance significance, defined as the ratio of the flight distance and its uncertainty, can be used as observable. The multiplicity of the tracks associated to a secondary vertex, the secondary vertex mass, assuming the pion mass for all tracks, or its energy are used.

Secondary-vertex candidates have to share less than 65% of their associated tracks with the primary vertex and the significance of the radial distance between both vertices has to be larger than 3σ . Candidates are rejected which have a radial distance of more than 2.5 cm with respect to the primary vertex, which have masses compatible with the K^0 mass or which exceed $6.5 \text{ GeV}/c^2$ to suppress vertices corresponding to decays of long-lived mesons or particle interactions with the detector material. The flight direction of each candidate has to be within a cone of $\Delta R < 0.5$ around the jet direction.

6.8.2. The Combined Secondary Vertex (CSV) algorithm

A commonly used algorithm for the identification of jets originating from b quarks is the Combined Secondary Vertex (CSV) algorithm. It combines the information of secondary vertices with the track-based information. This allows discrimination in cases where no secondary vertex is found and increases the efficiency of the CSV algorithm. Tracks with an impact parameter significance of $S_{IP} > 2$ can be combined in a "pseudo vertex", allowing for a subset of secondary vertex quantities.

The following variables are used in the CSV algorithm (the "no vertex" category only uses the last two variables):

- vertex category: real, "pseudo", or "no vertex"
- 2D flight distance significance: using the flight distance in the transverse plane
- secondary vertex mass
- number of tracks at the secondary vertex
- energy ratio carried by tracks at the secondary vertex with respect to all tracks in the jet
- pseudorapidities of the tracks at the secondary vertex with respect to the jet axis
- 2D IP significance of the first track that raises the invariant mass above the charm threshold of 1.5 GeV (tracks are ordered by decreasing IP significance and the mass of the system is recalculated after adding each track)
- number of tracks in the jet
- 3D IP significances for each track in the jet

Two likelihood ratios are build from these variables, one likelihood ratio to discriminate between b-quark and c-quark jets, and an other one to discriminate between b-quark and light-parton jets. These two likelihoods are needed, since charmed hadrons have similar properties as bottom hadrons. Using just the likelihood for light-flavour discrimination would lead to an higher fraction of c-quark jets which are identified as b-quark jet. As light-flavour quarks are the most important background for b-quark jet identification, the likelihood ratio for light-flavour discrimination is more important for b-quark jet identification than the other likelihood for c-quark jet discrimination. Both likelihoods are combined with prior weights of 0.25 and 0.75, respectively.

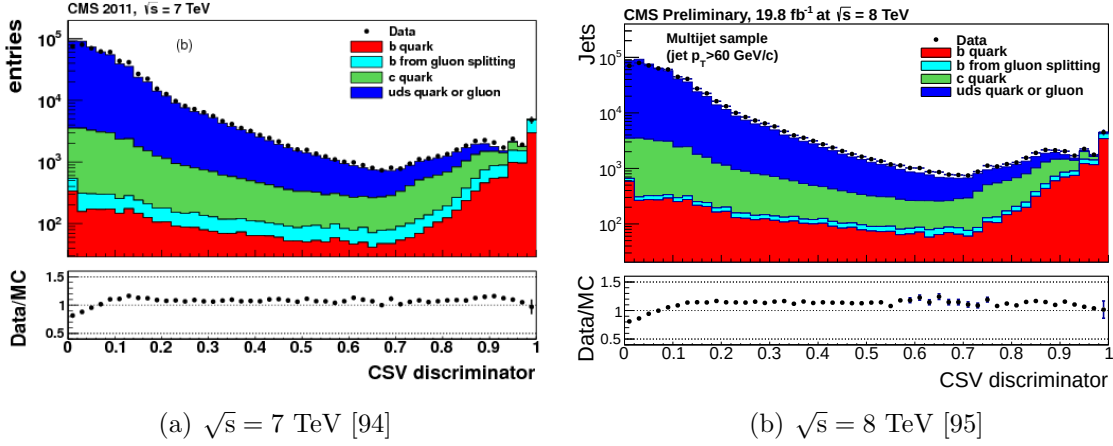


Figure 6.7.: Comparison of the observed CSV discriminator distribution with the MC background expectation after basic selection of QCD multijet events. The filled areas indicate the different flavour contributions for each bin, originating from b -quark jets (red), c -quark jets (green) and light-flavour jets (dark blue). Simulated events involving b quarks from gluon splitting (cyan) are indicated separately from other b -quark production-processes. The CSV distribution is shown for different run periods for LHC Run I (a) at $\sqrt{s} = 7$ TeV, and (b) at $\sqrt{s} = 8$ TeV.

The jet flavour can be identified for jets coming from event simulation. Generated hadrons are matched to the jet, and the heaviest parton within $\Delta R < 0.3$ to the jet axis defines the jet flavour. Figure 6.7 compares the observed CSV discriminator distribution with background expectation at different centre-of-mass energies from LHC Run I. Sample of inclusive multijet events were collected using single jet triggers with a p_T threshold of 60 GeV. QCD multijet events were generated with PYTHIA 6.4 using the Z tune [116] for $\sqrt{s} = 7$ TeV and the Z2* tune [114, 115] for $\sqrt{s} = 8$ TeV. Jets originating from b -quarks, c -quarks and light-flavour quarks are indicated in different colours. The number of simulated events of each component is normalised in such a way that the total number of simulated events fit the observed number of events. Jets originating from light-flavour tend to have lower CSV values, while b -quarks tend to have larger ones. Since charmed hadrons have similar properties to bottom hadrons, c -quark jets can have low and high CSV discriminator values. Due to the likelihood ratio to discriminate between b -quark and c -quark jets, some events have a low CSV discriminator value.

6.8.3. Efficiency of the Combined Secondary Vertex (CSV) algorithm

When selecting b -quark jets, it is crucial to know the efficiency of the used algorithm and the discriminator threshold. The misidentification efficiency is important, too. If a veto on b -quark jets is applied to suppress for example events originating from $t\bar{t}$ production, the same reasoning applies. These efficiencies are measured as function of jet p_T and η . They are calculated for observed and simulated events.

6. Event reconstruction

Identification efficiencies can be measured using QCD multijet events or $t\bar{t}$ events. In QCD multijet events, the b-quark jets have to be identified without using the algorithms for the efficiency measurement. The branching fraction of bottom hadrons which has a muon as decay product is significantly larger than for other hadrons and muons are efficiently identified in the CMS detector, therefore jets are selected for the b-quark identification measurement which contain a muon within $\Delta R = 0.4$ from the jet axis. As the top quark decays almost exclusively to a W boson and a b quark, there have to be at least two b-quark jets in $t\bar{t}$ events. Several methods exist for the efficiency and misidentification measurements [94]. Systematic uncertainties are coming from pileup, muon selection, jet-parton matching, PDFs, and trigger efficiencies. A systematic uncertainty is coming from the differences in the observed and simulated fraction of gluon splitting into $b\bar{b}$ pairs [94].

For the measurement of the misidentification efficiency, a sample enriched in light-flavour jets is used using inverted algorithms, which use tracks with negative IP values and secondary vertices with negative decay lengths. The efficiency for a jet not originating from a b-quark jet, $\varepsilon_{\text{data}}^-$, can be measured in collision data. A correction factor R_{light} taken from simulation corrects for differences in the efficiency between these inverted algorithms and the b-quark jet identification algorithms. It is defined as:

$$R_{\text{light}} = \varepsilon_{\text{MC}}^{\text{misid}} / \varepsilon_{\text{MC}}^-, \quad (6.15)$$

where $\varepsilon_{\text{MC}}^{\text{misid}}$ is the misidentification efficiency estimated from simulation. Then the misidentification efficiency in collision data can be evaluated as:

$$\varepsilon_{\text{data}}^{\text{misid}} = \varepsilon_{\text{data}}^- \cdot R_{\text{light}}. \quad (6.16)$$

Uncertainties are coming from fractions of c- and b-quark jets which are identified as jets originating from light-quark jets and the fraction of gluon jets, due to differences between light-flavour jets and jets from gluon splittings. Another uncertainty comes from long-lived particle decays, such as K_S^0 and Λ decays. Other sources for uncertainties are mismeasured tracks, IP sign-flips and pileup. A detailed description of the different efficiency and misidentification measurements and their uncertainties can be found in [94].

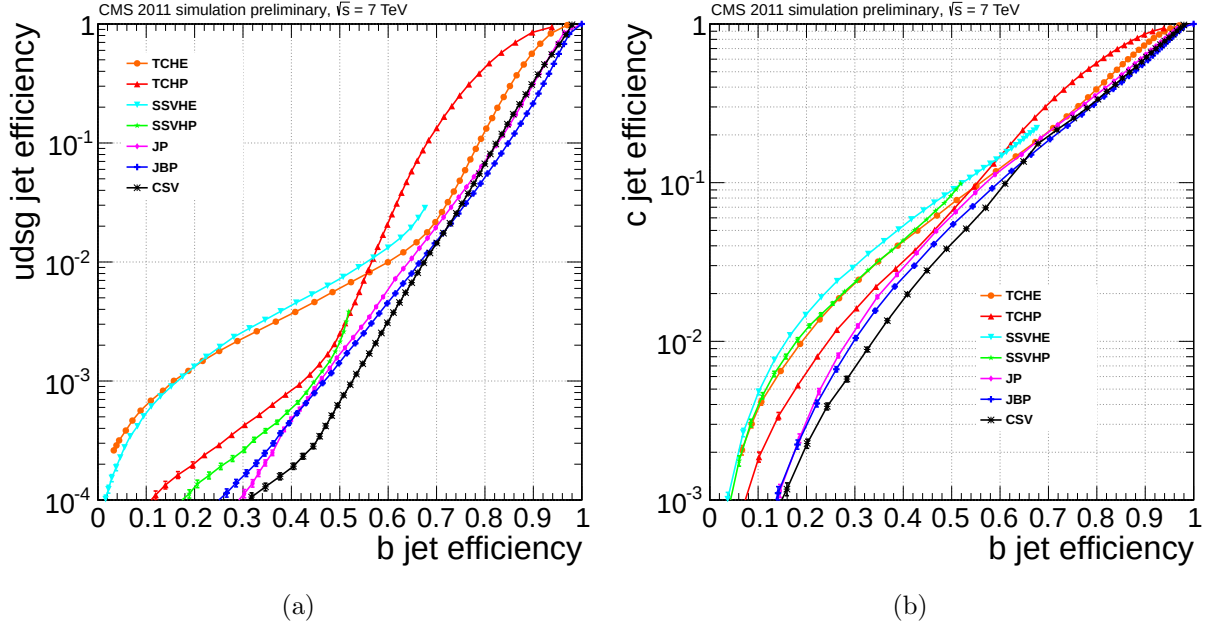


Figure 6.8.: Performance curves obtained from simulation for different b -quark jet identification algorithms. Jets with a momentum $p_T > 60$ GeV in a sample of QCD multijet events are used to obtain the efficiency values for identifying b -quark jets and the misidentification probability values for (a) light-flavour jets and (b) c -quark jets at $\sqrt{s} = 7$ TeV [94].

Figure 6.8 shows the performance curves for different b -quark jet identification algorithms including the CSV algorithm. The efficiency for identifying a b -quark jet and the misidentification probability of light-flavour or c -quark jets are shown for different operating points with a certain threshold on the discriminator value. The CSV algorithm has higher efficiencies for b -quark jets and lower misidentification probabilities than the other algorithms for most operating points. Three working points with a specified misidentification probability for light-flavour jets are defined. A loose working point “L” with a misidentification efficiency of 10%, a medium working point “M” with a misidentification efficiency of 1% for light-flavour jets and a tight working point “T” with a misidentification probability of 0.1%.

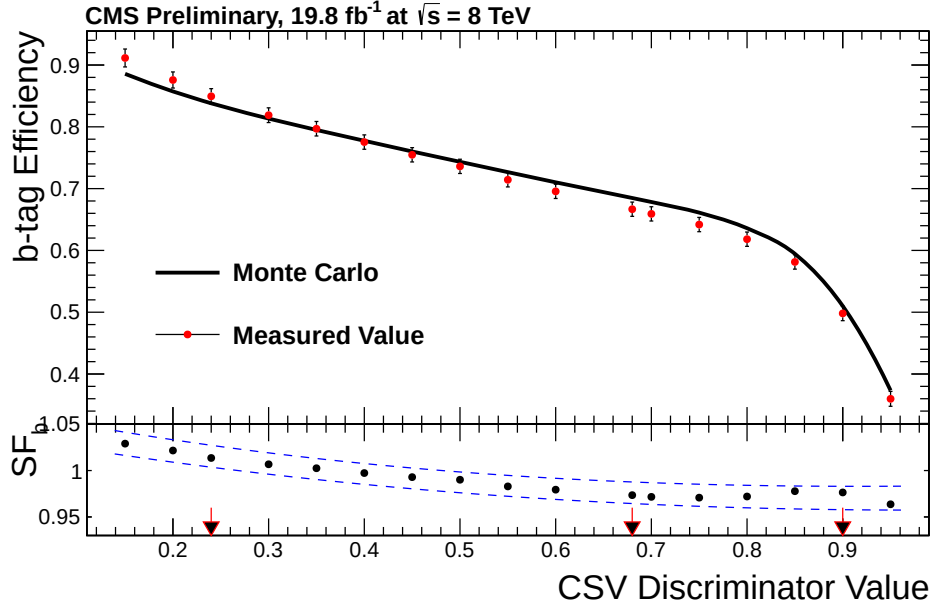


Figure 6.9.: Measured efficiency in observed events (red dots) and the predicted efficiency (black line) of the CSV algorithm to identify b-quark jets as function of the CSV discriminator threshold at $\sqrt{s} = 8$ TeV. Its corresponding scale factor SF_b is shown in the lower panel, where the blue dashed lines represent the combined statistical and systematic uncertainty. The arrows indicate the standard operating points, from left to right loose “L”, medium “M” and tight “T” [95].

The efficiencies measured from collision data and simulated events differ from each other (see Figure 6.9) due to hadronisation model uncertainties. The number of tracks depends on the decays of the b-hadrons. Also the number of gluon splittings and light-flavour jets depend on the used model and implementation of the parton shower. To correct for the differences in observed events and MC simulation, scale factors, which are defined as the ratio of the efficiency of observed events and the efficiency, can be applied for simulated jets. The MC information is used to identify the jet flavour. There are different scale factors for each flavour, SF_b for b-quark jets, SF_{light} for light-flavour jets, which depend on the jet momentum and pseudorapidity of the jet:

$$SF_b = \varepsilon_b^{\text{meas}} / \varepsilon_b^{\text{MC}}, \quad (6.17)$$

$$SF_{\text{light}} = \varepsilon_{\text{data}}^{\text{misid}} / \varepsilon_{\text{MC}}^{\text{misid}}. \quad (6.18)$$

Scale factors were derived for the $\sqrt{s} = 8$ TeV data set [95]. Multijet events were collected using single jet triggers with different thresholds. Simulated samples of QCD multijet events were generated with PYTHIA 6.4 using the Z2* tune [114, 115].

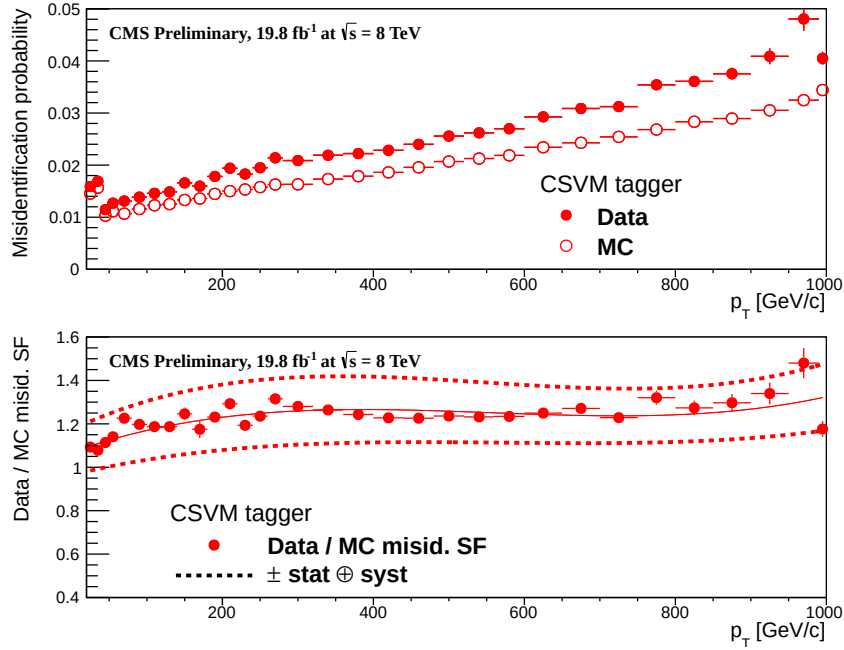


Figure 6.10.: Misidentification probabilities (top) and scale factors SF_{light} (bottom) of the CSV algorithm using the medium working point as function of jet momentum at $\sqrt{s} = 8$ TeV. The last bin includes all jets with $p_T > 1000$ GeV. Observed values are indicated by filled circles and values corresponding to simulated events are shown as open circles. The solid curve (bottom) is the result of a polynomial fit to the scale factors and the overall statistical and systematic uncertainties on the measurement are represented by dashed curves [95].

Figure 6.9 shows the scale factor SF_b as function of the CSV discriminator threshold including the working points for $\sqrt{s} = 8$ TeV. The misidentification probability in observed and simulated events as function of the jet momentum for $\sqrt{s} = 8$ TeV and the corresponding scale factors SF_{light} are shown in Figure 6.10.

Since it is difficult to distinguish c -quark jets from b -quark jets, it is also difficult to measure the scale factors for c -quark jets SF_c in collision data. Thus, b -quark jet scale factors are applied on c -quark jets in many analyses, assuming they are the same as for b -quark jets, with the corresponding uncertainty doubled [94].

6.8.4. Identification of b-quark jets in LHC Run II

Proton-proton collision data recorded by the CMS detector at a centre-of-mass energy of $\sqrt{s} = 13$ TeV corresponding to an integrated luminosity of about 2.6 fb^{-1} are analysed. The b-quark jet identification in LHC Run II [96] was modified compared to LHC Run I in order to cope with the new running conditions.

First, changes with regard to reconstructed objects are discussed. Then, a new algorithm for secondary-vertex reconstruction, the inclusive vertex finder algorithm [96, 117], is described. An optimised version of the CSV algorithm has been developed, referred to as CSVv2 algorithm. Finally, a new method for correcting simulated distribution to collision data regarding the efficiency of b-quark jet identification is discussed.

Reconstructed objects Jets are reconstructed by clustering PF candidates using the anti- k_T algorithm with a distance parameter $R = 0.4$, which is smaller compared to $R = 0.5$ used in LHC Run I. Jet flavour is determined by re-clustering the jet constituents including generator-level hadrons and partons. A jet flavour is assigned to a jet according to the clustered hadrons (partons) inside a jet. Higher priority is given to the b flavour, followed by the c flavour. If neither the b flavour nor the c flavour can be assigned to the jet, it is considered as light flavour.

Tracks and primary vertices are reconstructed and selected as in LHC Run I. There have been changes with respect to secondary vertex reconstruction. The secondary-vertex reconstruction algorithm used in LHC Run I is the adaptive vertex reconstruction (AVR) algorithm based on an adaptive vertex fit [101]. In LHC Run II, secondary vertices have to fulfil the same requirements as in LHC run I, but the flight direction of each secondary vertex has to be within a cone of $\Delta R < 0.4$ around the jet axis instead of $\Delta R < 0.5$ as in LHC Run I.

A new algorithm for secondary-vertex reconstruction has been implemented. The inclusive vertex finder (IVF) algorithm uses a different track selection and is not seeded from tracks associated to a reconstructed jet. This algorithm has a higher reconstruction efficiency of secondary vertices for b-quark jets (10%) and c-quark jets (15%) compared to the AVR algorithm. For light-flavour jets the efficiency to reconstruct a secondary vertex increases by 8%. Since secondary vertices are used in identification of b-quark jets, this has an influence on the misidentification rate. A better secondary vertex reconstruction can improve the discrimination between b-quark jets and other jets, but a higher reconstruction efficiency for secondary vertices could also lead to slightly more misidentified jets.

The Inclusive vertex finder (IVF) algorithm uses reconstructed tracks in the event as input. Selected tracks are required to have a transverse momentum exceeding 0.8 GeV and to have at least 8 hits in the silicon strips tracker. There is no requirement on the number of hits in the silicon pixel tracker. Their longitudinal impact parameter, d_z has to be smaller than 0.3 cm.

Displaced tracks are identified as seeds when their impact parameter, IP, has a value of at least $50 \mu\text{m}$ and a significance of at least 1.2σ . These seed tracks and nearby tracks are clustered based on their minimum distance relative to other tracks and the angles between them. Secondary vertex candidates are obtained by fitting the clusters with an adaptive vertex fitter. Candidates with a 2D (3D) flight distance significance below 2.5σ (0.5σ) are removed.

Vertices are also removed when they share more than 70% of their tracks with another vertex and the distance significance between them is less than 2σ . One of both such vertices is removed to avoid that same tracks appear in multiple vertices. Tracks still can be assigned to a primary and a secondary vertex, therefore tracks are discarded from the secondary vertex when they have less than 1 hit in the pixel tracker or when they are more compatible with the primary vertex. The secondary vertex is refitted if at least 2 tracks are associated to the secondary vertex candidate. If it shares 20% of its tracks with another candidate and the distance significance between both candidates is less than 10σ , the vertex is removed.

Secondary vertices are selected with tighter requirements as the vertices reconstructed with the AVR algorithm. The fraction of shared tracks between the candidate and the primary vertex has to be smaller than 79%, instead of 65%. Their flight direction has to be within a cone of $\Delta R < 0.3$ with respect to the jet axis, compared to $\Delta R < 0.4$ for AVR secondary-vertices. The flight-distance significance of the secondary vertices is required to be larger than 2σ , instead of 3σ .

Combined secondary vertex (CSVv2) algorithm is an optimised version of the CSV algorithm used in LHC Run I. It combines information coming from displaced tracks and secondary vertices by using a multivariate technique. The main differences with LHC Run I are the different secondary-vertex algorithm used, the input variables and their combination. There are two variants of the CSVv2 algorithm, depending whether IVF or AVR secondary-vertices are used.

6. Event reconstruction

Variables which have been added are:

- number of secondary vertices
- angle between secondary vertex and jet axis
- ratio of the transverse momentum of summed track four-momenta and the jet momentum
- decay length and the angle between the track and the jet

As in LHC Run I, there are 3 independent vertex categories, real, “pseudo”, and “no vertex”, but the way how the input variables are combined is different. They were combined with a likelihood ratio in LHC Run I. In the new version, a discriminator is build for each vertex category by combining information of discriminating variables using a multilayer perceptron² with one hidden layer. A likelihood ratio is build from these three discriminators taking the fraction of each jet flavour into account.

Figure 6.11 shows the observed and simulated CSV discriminator distributions for LHC Run I and Run II. Multijet samples were collected using jet triggers with different thresholds on the jet transverse momentum. QCD multijet events were generated with PYTHIA 6.4 using the Z tune [116] for events at $\sqrt{s} = 7$ TeV and the Z2* tune [114,115] for events at $\sqrt{s} = 8$ TeV. For LHC Run II, multijet events were generated with PYTHIA 8 for a centre-of-mass energy of $\sqrt{s} = 13$ TeV using the CMS Underlying Event Tune for PYTHIA 8 (CUETP8M1) [118]. Top quark pair production as well as single top quark production in association with a W boson were generated using POWHEG and PYTHIA 8 with the CUETP8M1 tune. A dilepton sample was collected at $\sqrt{s} = 13$ TeV by selecting events with at least one isolated electron and at least one isolated muon. Events are further considered if both leptons have a transverse momentum exceeding 20 GeV and if they contain at least 2 jets with a transverse momentum exceeding 20 GeV. In Figure 6.11, the number of simulated events of each component is normalised in such a way that the total number of simulated events fit the observed number of events. The distribution of the CSVv2 algorithm differs from the CSV distribution of LHC Run I. There are also dependencies on the selected sample as shown in Figure 6.11(c) and (d). Shape and fraction of the different components differ, but the main features stay the same. The low discriminator-value region is dominated by light-flavour jets and the high discriminator-value region is dominated by b-quark jets. Jets originating from c-quark jets peak contribute to the full range of the discriminator distribution. The agreement between MC simulation and collision data gets worse for low discriminator values, but it is very good elsewhere.

²Artificial neural network with multiple layers of nodes that maps input data onto a set of outputs [105].

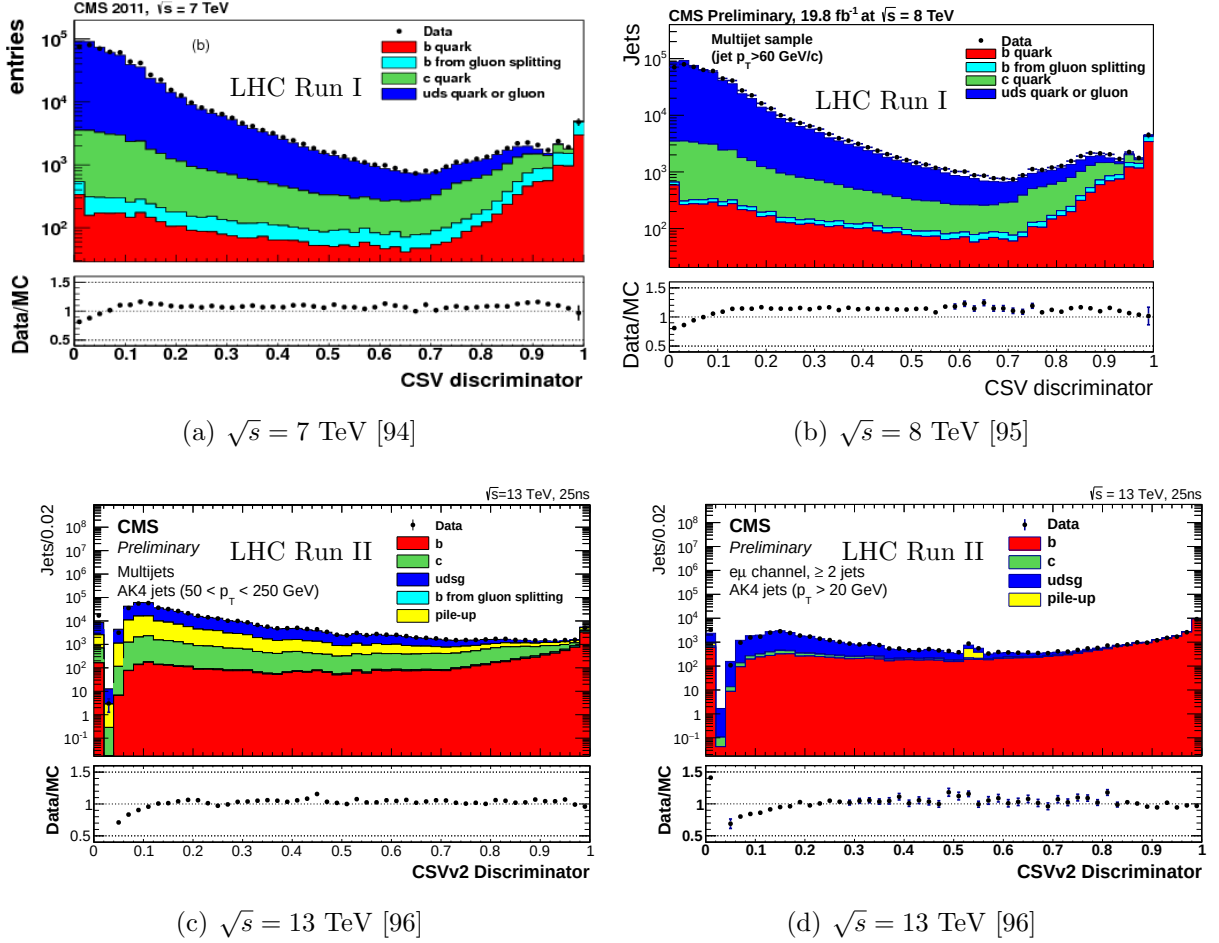


Figure 6.11.: Comparison of the observed CSV discriminator distribution with the background expectation after basic selection of QCD multijet events (a, b, c) and after basic selection of dilepton $t\bar{t}$ events (d). Distribution is shown for different running periods at $\sqrt{s} = 7$ TeV (a), $\sqrt{s} = 8$ TeV (b) and $\sqrt{s} = 13$ TeV (c, d). The filled areas indicate the different flavour contributions for each bin, originating from b -quark jets (red), c -quark jets (green) and light-flavour jets (dark blue). Simulated events involving b quarks from gluon splitting (cyan) are indicated separately from other b -quark production-processes. There are contributions of jets originating from pileup (yellow) for LHC Run II.

Identification efficiency from $t\bar{t}$ events

Same techniques as for LHC Run I are used for measuring the identification efficiencies in collision data and for estimating them in MC simulation. Distributions using generated events have to be corrected to match events observed in collision data. A new reweighting technique [96,119], Reweight method, has been developed which calibrates the full discriminator shape instead of correcting the identification efficiency after requiring the events to pass a specific working point. Scale factors are derived as function of the discriminator value, jet p_T and jet η in order to scale identification and misidentification efficiencies for b-quark jet identification in MC simulation to the one observed in collision data.

Dilepton events with two isolated leptons are selected and a tag and probe method is used to measure scale factors for b-quark jets and light-flavour jets simultaneously. Top events from $t\bar{t}$, tW and $\bar{t}W$ production were generated using POWHEG and PYTHIA 8 with the CUETP8M1 tune. The two isolated leptons are required to have an opposite sign, a transverse momentum exceeding 20 GeV, and an invariant mass M_{ll} larger than 12 GeV. There have to be exactly two jets with a transverse momentum exceeding 30 GeV. Events with same flavour leptons have to fulfil some extra requirements to reduce the background from Z+jets events. A veto around the Z boson mass is applied, i.e. $|M_{ll} - M_Z| > 10$ GeV. The missing transverse momentum, $\cancel{E}_T$, is required to be larger than 30 GeV.

Scale factors for b-quark jets, SF_b , are measured by requiring the tag jet to pass the medium working point of the b-quark jet identification algorithm. Contributions from non b-quark jets are subtracted using simulated events. When measuring scale factors for light-flavour jets, the event selection is enriched in Z+jets events by requiring two leptons with same flavour and inverting the veto on the Z boson mass and $\cancel{E}_T$. A b-quark jet veto is applied on the tagged jet using the loose working point of the b-quark jet identification algorithm. Expected contributions from b-quark and c-quark jets are subtracted. Since scale factors of the other contributions have an impact on measured scale factors, an iterative procedure is applied. No scale factors are applied in the first iteration. In the second iteration, scale factors obtained in the previous iteration are used. The scale factor for c-quark jets is set to unity assuming the uncertainty to be twice as large as the one for the scale factor of b-quark jets. When the scale factors obtained in the current iteration are stable with respect to the scale factors of the previous iteration, the iterative procedure stops. Systematic uncertainties include the jet energy scale uncertainty (JES), contamination of jets with a different flavour which is varied by 20%, the uncertainty of c-quark jet scale factors, and the statistics of the simulated samples.

This method is useful for analyses which apply a veto on b-tagged jets because scale factors are derived for the identification of b-quark jets and not for vetoing them.

7. Impact of heavy flavour modelling on the misidentification of c-quark jets

Since the properties of charmed hadrons are very similar to those of bottom hadrons, jets originating from c-quark jets are the main background for the identification of b-quark jets. Scale factors are derived for correcting differences between measured efficiencies in collision data and simulation. As the properties of charmed and bottom hadrons are very similar, it is difficult to derive the scale factors for c-quark jets. Scale factors for b-quark jets were derived at $\sqrt{s} = 8$ TeV using PYTHIA 6.4 and the Z2* tune [114, 115]. As described in the previous chapter in Section 5, different MC generators use different models for hadronisation and hadron decays. Since b-quark jet identification algorithms rely on the differences between bottom, charmed and other hadrons, differences in the modelling of hadronisation and of the charmed hadrons can affect the misidentification probability, therefore a precise simulation of these quantities is important.

In this chapter, differences in heavy flavour modelling between PYTHIA 6 and HERWIG 6 and their effect on the c-quark jet misidentification probability are studied for proton-proton collisions at $\sqrt{s} = 8$ TeV. One of the main results of this study is that the misidentification probability depends on the type of the charmed hadron which is produced during the hadronisation. First, the simulation, selection and categorisation of jets used in the study are described. The dependence of the misidentification efficiency on the charmed hadron type is presented. Then, it is discussed which differences in heavy flavour modelling could affect the misidentification probability of c-quark jets. Differences in the heavy flavour modelling of charm quarks between PYTHIA 6 and HERWIG 6 are presented. Afterwards, the effect on the misidentification probability is investigated and the accuracy of the heavy flavour modelling is studied by comparison to measurements. Finally, systematic uncertainties on the misidentification probability of c-quark jets as b-quark jets are discussed and it is studied if the systematic uncertainties cover the differences in heavy flavour modelling.

7.1. Simulation, selection and categorisation of jets

Long-lived charmed hadrons have similar properties as long-lived bottom hadrons, but a less hard fragmentation, a shorter lifetime and smaller masses. Table 7.1 lists the most commonly produced charmed mesons and baryons, their quark content, their mass and their life time. Charmed mesons which contain only light-flavour quarks, $D^\pm$ and $D^0/\bar{D}^0$, are dominantly produced. In the following, the remaining charmed hadrons of Table 7.1 are referred to as *other charmed hadrons*.

Table 7.1.: Charmed hadrons and their quark content [120].

	Charmed hadrons	Quark content	Mass [MeV]	$c\tau$ [μm]
Charmed mesons	D^+/D^-	$c\bar{d}/\bar{c}d$	1869.62 ± 0.15	311.8
	$D^0/\bar{D}^0$	$c\bar{u}/\bar{c}u$	1864.86 ± 0.13	122.9
Charmed strange mesons	D_s^+/D_s^-	$c\bar{s}/\bar{c}s$	1968.49 ± 0.32	149.9
Charmed baryons	Λ_c^+/Λ_c^-	$udc/\bar{u}d\bar{c}$	2286.46 ± 0.14	59.9
Charmed strange baryons	Ξ_c^0/Ξ_c^-	$dsc/\bar{d}s\bar{c}$	$2470.88^{+0.34}_{-0.80}$	33.6
	Ξ_c^+/Ξ_c^-	$usc/\bar{u}s\bar{c}$	$2467.8^{+0.4}_{-0.6}$	132
	$\Omega_c^0/\bar{\Omega}_c^0$	$ssc/\bar{s}s\bar{c}$	2695.2 ± 1.7	21

QCD multijet events generated with PYTHIA 6 and HERWIG 6 are used to study the effect of the heavy flavour modelling on the b -quark jet identification for proton-proton collisions at $\sqrt{s} = 8$ TeV. The detector response is modelled using GEANT4 [85]. PYTHIA parameters are set to the Z2* tune [114, 115], while HERWIG events are generated using the standard HERWIG 6 tune. A list of these samples can be found in Table 7.2. They are binned in $\hat{p}_T$, where $\hat{p}_T$ is the transverse momentum of the outgoing partons in the parton-parton centre-of-mass frame for 2-to-2 parton-parton scattering.

Table 7.2.: Simulated QCD multijet samples binned in $\hat{p}_T$ for proton-proton collisions at $\sqrt{s} = 8$ GeV [121]. PYTHIA 6 events are generated using the Z2* tune, while for the HERWIG 6 events the standard CMS HERWIG tune is used.

$\hat{p}_T$	PYTHIA 6	HERWIG 6
30-50	QCD_Pt-30to50_TuneZ2star_8TeV_pythia6	QCD_Pt_30to50_CTEQ6L1_8TeV_herwig6
50-80	QCD_Pt-50to80_TuneZ2star_8TeV_pythia6	QCD_Pt_50to80_CTEQ6L1_8TeV_herwig6
80-120	QCD_Pt-80to120_TuneZ2star_8TeV_pythia6	QCD_Pt_80to120_CTEQ6L1_8TeV_herwig6
120-170	QCD_Pt-120to170_TuneZ2star_8TeV_pythia6	QCD_Pt_120to170_CTEQ6L1_8TeV_herwig6
170-300	QCD_Pt-170to300_TuneZ2star_8TeV_pythia6	QCD_Pt_170to300_CTEQ6L1_8TeV_herwig6
300-470	QCD_Pt-300to470_TuneZ2star_8TeV_pythia6	QCD_Pt_300to470_CTEQ6L1_8TeV_herwig6
470-600	QCD_Pt-470to600_TuneZ2star_8TeV_pythia6	QCD_Pt_470to600_CTEQ6L1_8TeV_herwig6
600-800	QCD_Pt-600to800_TuneZ2star_8TeV_pythia6	QCD_Pt_600to800_CTEQ6L1_8TeV_herwig6
800-1000	QCD_Pt-800to1000_TuneZ2star_8TeV_pythia6	QCD_Pt_800to1000_CTEQ6L1_8TeV_herwig6
>1000	QCD_Pt-1000to1400_TuneZ2star_8TeV_pythia6	QCD_Pt_1000_CTEQ6L1_8TeV_herwig6

The reconstruction of jets is described in Section 6.6 and the algorithms for the identification of b-quark jets and the assignment are discussed in Section 6.8. Jet flavour and how a jet flavour is assigned to a jet is also introduced in both sections.

Reconstructed jets with a momentum $p_T > 20$ GeV and $|\eta| < 2.5$ are selected, which were assigned to be c-quark jets. Jets are divided into categories according to types of charmed hadrons. The reconstructed c-quark jets are matched to charmed hadrons, which do not decay further into other charmed hadrons. The type of the charmed hadron is assigned to the jet, which is in a cone of $\Delta R = 0.4$ around the jet axis and has the smallest ΔR relative to the jet axis. This angular distance ΔR has been chosen instead of the jet cone to ensure that tracks of charmed-hadron decay-products are assigned to the jet. Three categories of c-quark jets are defined, c-quark jets with a matched $D^\pm$ meson, c-quark jets with a matched $D^0/\bar{D}^0$ meson, and all other c-quark jets with a matched charmed hadron.

Figure 7.1 shows the performance of the CSV algorithm in dependence of the charmed hadron type, characterised by the corresponding b jet efficiencies and c jet misidentification probabilities. The misidentification probability of c-quark jets depends on the charmed hadron type which is produced during hadronisation. When D^+/D^- mesons are produced, the misidentification probability is higher than for the other charmed hadrons in general, since their decay length is larger. The category *other charmed hadrons* has the lowest misidentification probability. Their production rate in hadronisation is smaller, because of their quark content the formation of these hadrons is less probable. The likelihood ratio for discriminating b-quark jets from c-quark jets is optimised for the dominant produced charmed hadrons, which have other properties. This could explain the lower misidentification probability.

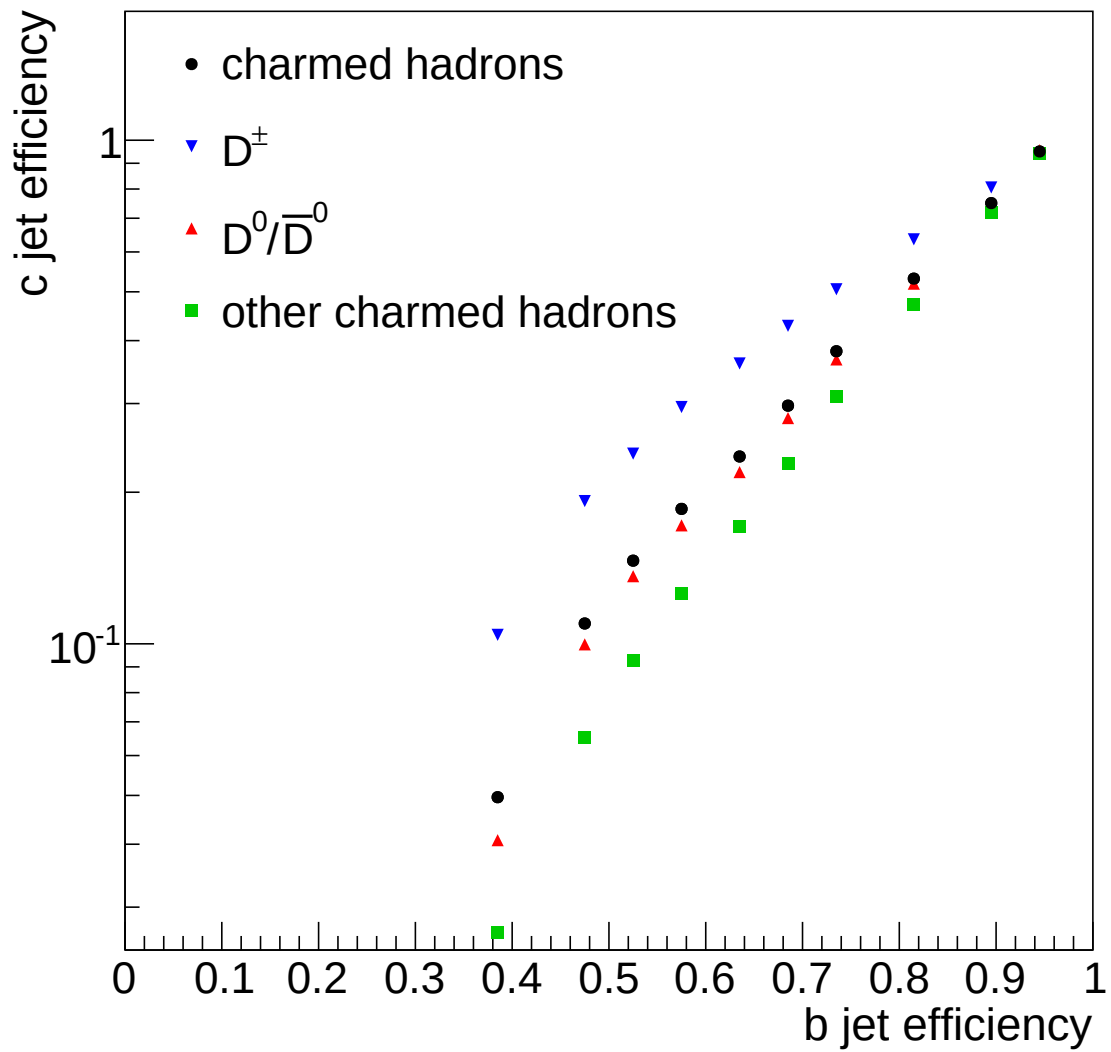


Figure 7.1.: Performance of the CSV algorithm in dependence on the charmed hadron type. Jets with a momentum $p_T > 20$ GeV in a MC sample of QCD multijet events are used to estimate the efficiency values for b -quark jets and the misidentification probability values for c -quark jets which contain a charmed hadron.

7.2. Differences in heavy flavour modelling

As the performance of the CSV algorithm depends on the charmed hadron type and different models are used by PYTHIA and HERWIG for hadronisation, differences in heavy-flavour modelling could also affect the c-quark jet misidentification probabilities. Tracks are used by the CSV algorithm to identify the bottom hadron decays, therefore differences in decay modelling of charmed hadrons could lead to different misidentification probabilities. Track reconstruction depends on the track momentum. The higher the hadron energy fraction is the higher is the momentum of the hadron decay products which lead to tracks in the detector used by the CSV algorithm. Different fragmentation modelling can thus affect the misidentification probability.

These three possible differences in heavy flavour modelling could affect the CSV misidentification probability of c-quark jets. First, it is investigated if there are differences in the charmed hadron production fractions between PYTHIA 6 and HERWIG 6. Then, differences in the decay modelling of charmed hadrons are studied. Next, the differences of model the fragmentation are studied. Finally, it is examined if these differences affect the CSV misidentification probability.

7.2.1. Charmed hadron production fractions

Charmed hadrons have slightly different masses leading to a different secondary vertex mass and a different momentum distribution of its decays products. Different decay chains are involved, thus the track multiplicity also depends on the charmed hadron type. This leads to a dependence of the performance of the CSV algorithm on the charmed hadron type, which is not negligible, as it can be seen in Figure 7.1.

The production fraction of a charmed hadron can be estimated by counting the c-quark jets with this type and divide this number by the total number of c-quark jets which contain a charmed hadron. Table 7.3 lists the fraction of the charmed hadron type to the total number of c-quark jets containing a charmed hadron for PYTHIA 6 and HERWIG 6. The ratio of these fractions is given for each hadron type, too.

7. Impact of heavy flavour modelling on the misidentification of c-quark jets

Table 7.3.: Charmed hadron type fractions in PYTHIA 6 and HERWIG 6 and their ratio. Simulated multijet events were generated with PYTHIA 6 and HERWIG 6 for proton-proton collisions at $\sqrt{s} = 8$ TeV. The charmed hadron type fraction is calculated by counting the c-quark jets with this hadron type and divide by the number of c-quark jets with a charmed hadron. Statistical uncertainties are negligible.

Type	PYTHIA	HERWIG	PYTHIA/HERWIG
D^+/D^-	18.39%	27.01%	0.6809
$D^0/\bar{D}^0$	61.5%	58.4%	1.0526
Other charmed hadrons	20.13%	14.58%	1.3805

More than half of the c-quark jets with a charmed hadron contain a D^0 or $\bar{D}^0$ hadron. Multijet samples generated with PYTHIA 6 have similar amounts of c-quark jets containing a D^+/D^- hadron or a hadron of the other charmed hadron category. In HERWIG 6, twice as many c-quark jets contain a D^+/D^- hadron as c-quark jets containing a hadron of the *other charmed hadrons* category.

As D^+/D^- hadrons have a higher misidentification probability than other charmed hadrons, this difference in charmed hadron production fractions will lead to a difference in the misidentification probability between PYTHIA and HERWIG as well.

In conclusion, there are large deviations in the charmed hadron production fractions between PYTHIA 6 and HERWIG 6.

7.2.2. Charmed hadron decays

The CSV algorithm uses the track-based life time information of the hadrons, such as the number of tracks in a jet and the number of tracks at the secondary vertex. Thus, the misidentification probability depends on the track multiplicity which is related to the hadron decay. In the case of the c-quark jet misidentification probability, the track multiplicity is related to the charged-daughter multiplicity of the charmed hadron. Generator information is used to determine the charged-daughter multiplicity. As decay tables are parameters of the event generators, the charged-daughter multiplicity can differ between PYTHIA 6 and HERWIG 6.

Figure 7.2 shows the charged-daughter multiplicities for each hadron type category and all c-quark jets with a charmed hadron. There are differences in the charged-daughter multiplicity between PYTHIA and HERWIG, especially for higher daughter multiplicities. Large deviations are in the other charmed hadron category. These differences originate from differences in the charmed hadron production fractions and could also come from differences in the decay modelling.

In summary, there are small differences in the modelling.

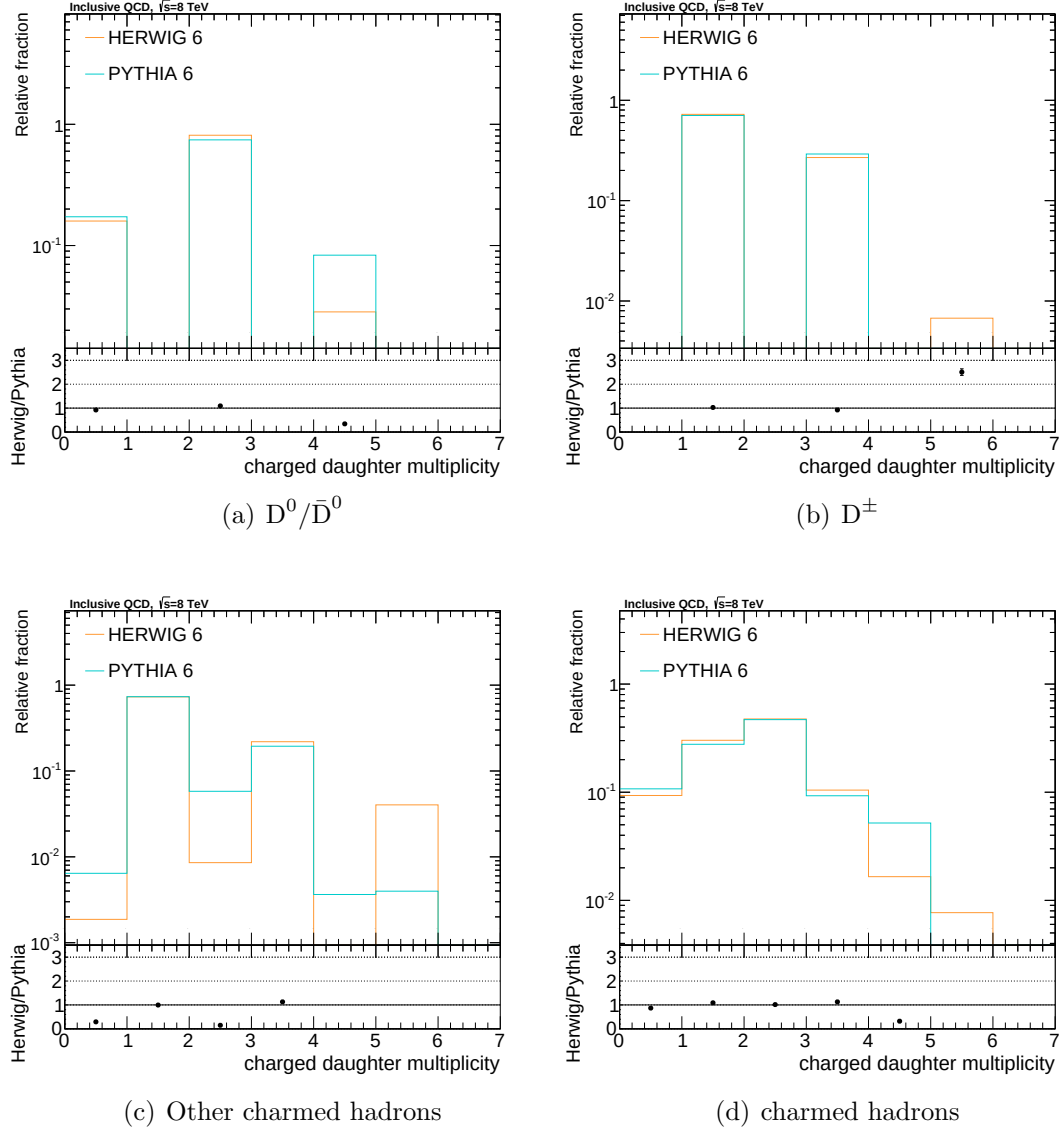


Figure 7.2.: Charged-daughter multiplicity in PYTHIA 6 and HERWIG 6 for (a) c-quark jets containing a D^0 or $\bar{D}^0$ meson, (b) for c-quark jets with a D^+ or D^- meson, for (c) for c-quark jets with a charmed hadron of the other charmed hadron category, and for (d) c-quark jets with a charmed hadron. Multijet samples were generated for proton-proton collisions at a centre-of-mass energy of $\sqrt{s} = 8$ TeV.

7.2.3. Charmed hadron momentum fraction

The CSV discriminator depends on the jet momentum and on the track information, which is also momentum dependent. Also the impact parameters of tracks assigned to the secondary vertex and their resolution depend on the track momentum, and are correlated to the charmed-hadron momentum, which in turn depends on the jet momentum.

The charmed hadron momentum fraction p_T^D/p_T^J is a better observable than the jet momentum p_T^J and the charmed hadron momentum p_T^D by itself. It takes correlation of both observables into account and it is easier to compare simulated distributions of p_T^D/p_T^J to measured ones obtained at experiments for charmed hadron physics with a lower centre-of-mass energy.

Figure 7.3 shows the momentum fraction in dependence on the charmed hadron category. They are approximately the same at lower momentum fractions but differ at higher momentum fractions. D^0 or $\bar{D}^0$ mesons carry a very high fraction of the jet momentum more often than other charmed hadrons.

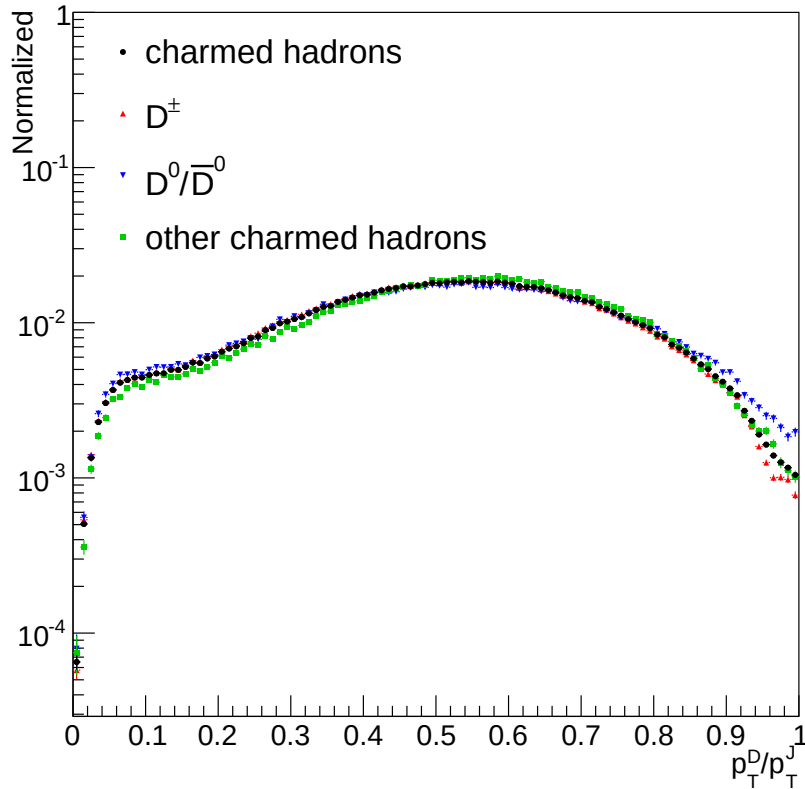


Figure 7.3.: Charmed hadron momentum fraction on the jet momentum for all charmed hadrons and the different charmed hadron types. QCD multijet events were generated with PYTHIA 6 for proton-proton collisions at a centre-of-mass energy of $\sqrt{s} = 8$ TeV.

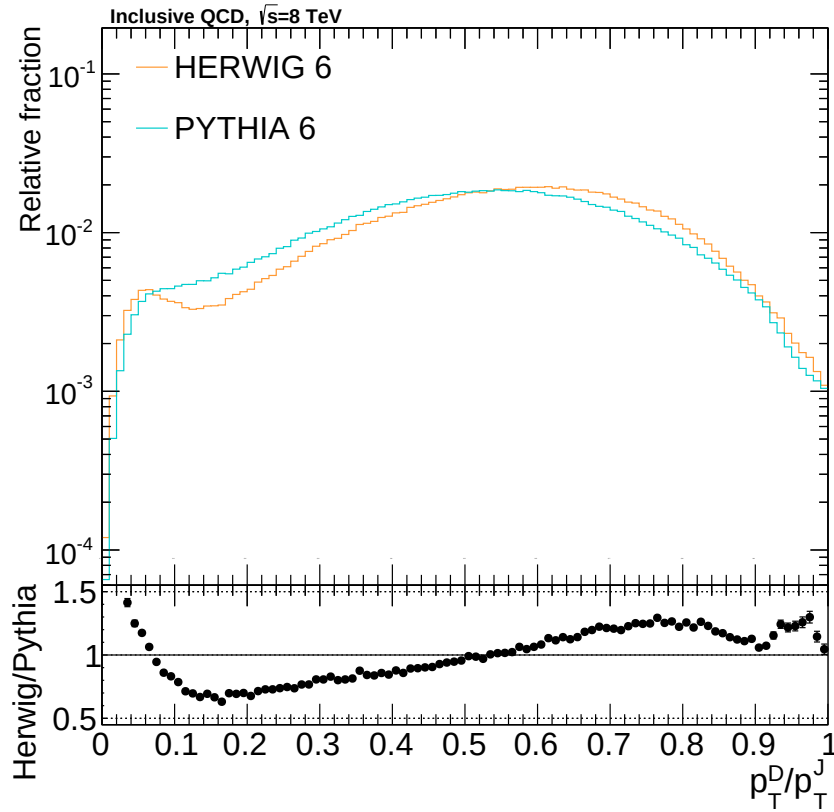


Figure 7.4.: Charmed hadron momentum fraction in PYTHIA 6 and HERWIG 6 for c-quark jets with a charmed hadron. Multijet QCD events were generated for proton-proton collisions with a centre-of-mass energy of $\sqrt{s} = 8$ TeV.

Figure 7.4 shows the charmed hadron momentum fraction in PYTHIA 6 and HERWIG 6 for c-quark jets. It is higher on average for HERWIG 6 for PYTHIA 6. A small part of this difference could come from the different charmed hadron production fractions, but $D^0/\bar{D}^0$ mesons carry a higher fraction of the jet momentum and the HERWIG sample has a smaller fraction of these mesons.

PYTHIA and HERWIG use different hadronisation models, therefore some differences in the momentum fraction distributions are expected.

In summary, charmed hadrons generated with PYTHIA 6 carry a different momentum fraction than charmed hadron which are generated with HERWIG 6.

7.2.4. Misidentification probability of c-quark jets

It is important to examine if the differences between PYTHIA 6 and HERWIG 6 in the heavy flavour modelling of charmed hadrons also affect the CSV discriminator and the performance of the CSV algorithm. Figure 7.5 shows the CSV discriminator and the performance of the CSV algorithm for these MC generators. The CSV discriminator distribution is shifted to higher values, when the events were generated with HERWIG 6. This leads to a higher misidentification probability for c-quark jets generated with HERWIG 6. Table 7.4 lists the misidentification probability for c-quark jets with a charmed hadron at each CSV working point. The amount of c-quark jet background depends on the selection on the CSV discriminator. The higher its value the less background is in the remaining sample of events. The largest deviations are for c-quark jets with high CSV discriminator values, therefore a tighter selection on the CSV discriminator leads to larger relative differences D_{rel} and systematic uncertainty in the misidentification probability of c-quark jets. The relative difference in the misidentification probability is defined as:

$$D_{rel} = \frac{\varepsilon_{\text{PYTHIA}} - \varepsilon_{\text{HERWIG}}}{\varepsilon_{\text{PYTHIA}}}, \quad (7.1)$$

where $\varepsilon_{\text{PYTHIA}}$ and $\varepsilon_{\text{HERWIG}}$ are the c-quark jet misidentification probabilities for PYTHIA 6 and HERWIG 6, respectively. This difference is in the order of 20% for the medium working point.

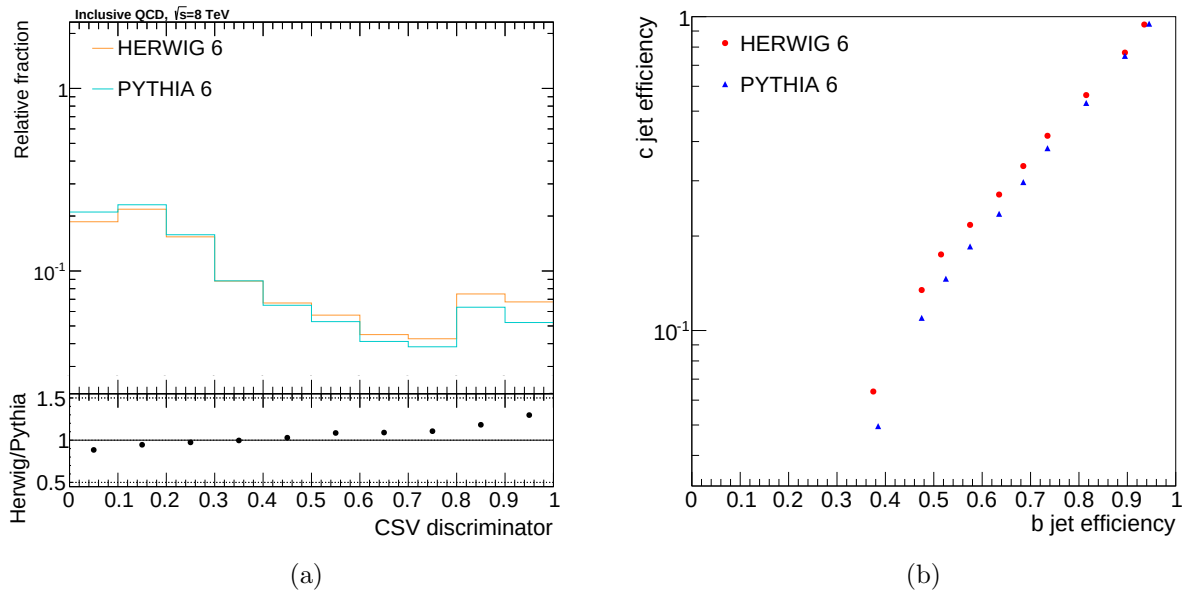


Figure 7.5.: CSV discriminator (a) and performance of the CSV algorithm (b) for c-quark jets with a charmed hadron in PYTHIA 6 and HERWIG 6.

Table 7.4.: Misidentification probability of c-quark jets with a charmed hadron for each CSV working point in PYTHIA 6 and HERWIG 6. The relative difference is defined in equation 7.1 QCD multijet events were generated for $\sqrt{s} = 8$ TeV.

Working point	CSV discriminator	PYTHIA 6	HERWIG 6	Relative difference
Loose	> 0.244	44.3%	48.1%	8.6%
Medium	> 0.679	15.0%	18.1%	20.8%
Tight	> 0.898	5.0%	6.5%	30.9%

7.2.5. Conclusions

Identification of b-quark jets is important for many analyses. It is crucial to know the efficiency and misidentification probabilities for each b-quark jet identification algorithm. Scale factors are derived for correcting deviations in observed and simulated efficiencies. This is done using simulated QCD multijet events which were generated with PYTHIA. However, PYTHIA and HERWIG use different models for hadronisation and hadron decay. This has an effect on the charmed hadron production fraction, the charged hadron daughter multiplicity, and the charged hadron momentum fraction. As could be shown, this leads to a difference in the performance of the CSV algorithm between c-quark jets which were generated with PYTHIA 6 and c-quark jets which were generated using HERWIG 6. This difference is in the order of 20% for the medium working point.

7.3. Reweighting HERWIG to PYTHIA values

The misidentification probability of c-quark jets generated with HERWIG 6 is higher than for PYTHIA 6. It is important to know, where the differences in the performance of the CSV algorithm for c-quark jets come from. This is done by reweighting c-quark jet distributions generated with HERWIG 6 to the one predicted by PYTHIA 6.

First, a single possible source is reweighted and its effects on the CSV discriminator distribution is discussed. Then, charmed hadron production fractions, charged-daughter multiplicities, and the charmed hadron momentum fractions are reweighted together to investigate if there are other sources for the difference in the performance of the CSV algorithm for c-quark jets between PYTHIA and HERWIG.

7.3.1. Reweighting one quantity

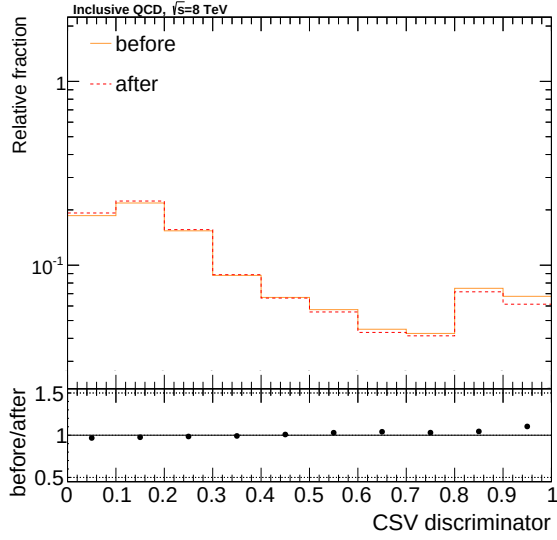
Charmed hadron production fractions, charged-daughter multiplicity of charmed hadrons and the charmed hadron momentum fractions are reweighted, separately, to study if the difference in this quantity is partly responsible for the difference in the misidentification probability of c-quark jets between PYTHIA 6 and HERWIG 6.

This is done by reweighting c-quark jets generated with HERWIG 6 in such a way that the distribution in the reweighted quantity match the PYTHIA 6 distribution. Quantities which depend on the charmed hadron category, like the charged daughter multiplicity and the momentum fraction, are reweighted for each hadron category. Then, it is checked how the distribution of the CSV discriminator is affected and if there are still remaining differences between HERWIG 6 and PYTHIA 6.

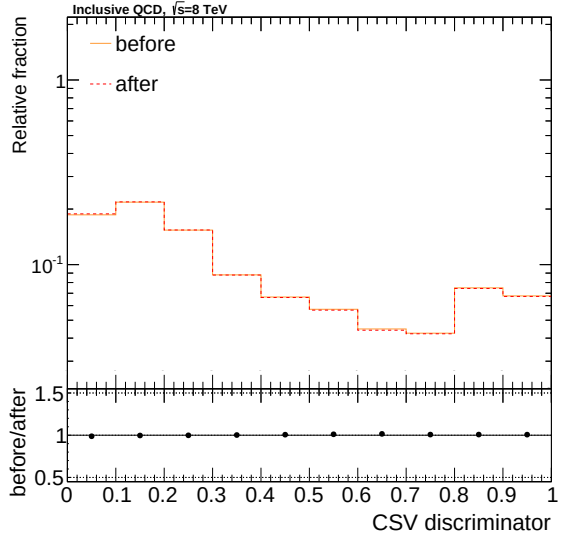
Figure 7.6(a) shows the CSV discriminator distribution of c-quark jets before and after reweighting c-quark jets generated by HERWIG 6 to predicted PYTHIA 6 production fractions. Jets with a high discriminator value are mostly affected.

In Figure 7.6(b), differences in the charmed hadron decay modelling are corrected by reweighting the charged-daughter multiplicity for each charmed hadron category. The CSV distribution is hardly affected by this.

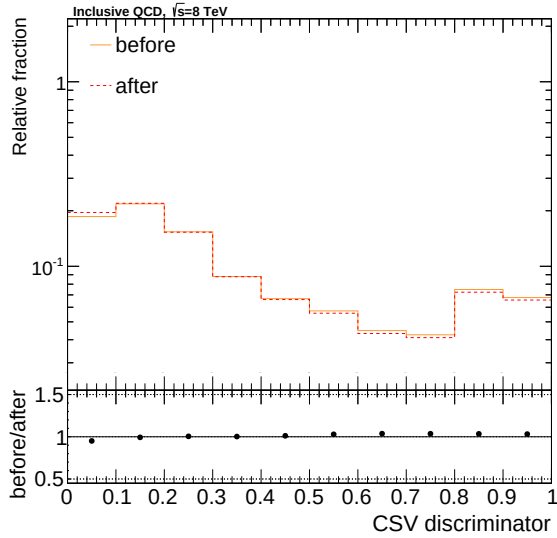
Charmed hadron momentum fractions depend on the hadron type, therefore they are reweighted for each type. Weights are calculated with a binning of 0.1 in the momentum fraction. Figure 7.6(c) shows, c-quark jets with higher CSV discriminator values are affected.



(a) reweight of production fractions



(b) reweight of decays



(c) reweight of momentum fraction

Figure 7.6.: CSV discriminator distribution of c-quark jets before and after reweighting HERWIG 6 (a) charmed hadron production fractions, (b) charged-daughter multiplicity of charmed hadrons, or (c) momentum fractions of charmed hadrons to PYTHIA 6 predictions. Multijet c-quark jets were generated for a centre-of-mass energy of $\sqrt{s} = 8$ TeV proton-proton collisions.

7. Impact of heavy flavour modelling on the misidentification of c -quark jets

Figure 7.7 shows the CSV discriminator distribution of c -quark jets with a charmed hadron generated with PYTHIA 6 and HERWIG 6, respectively. Distributions are shown before reweighting c -quark jets generated with HERWIG 6 to PYTHIA 6 predictions. Reweighting either charmed hadron production fractions or reweighting charmed hadron momentum fraction alone, does not remove all the differences between PYTHIA and HERWIG in the CSV discriminator distribution. Thus, either both or another quantity are responsible for the differences in the discriminator distribution.

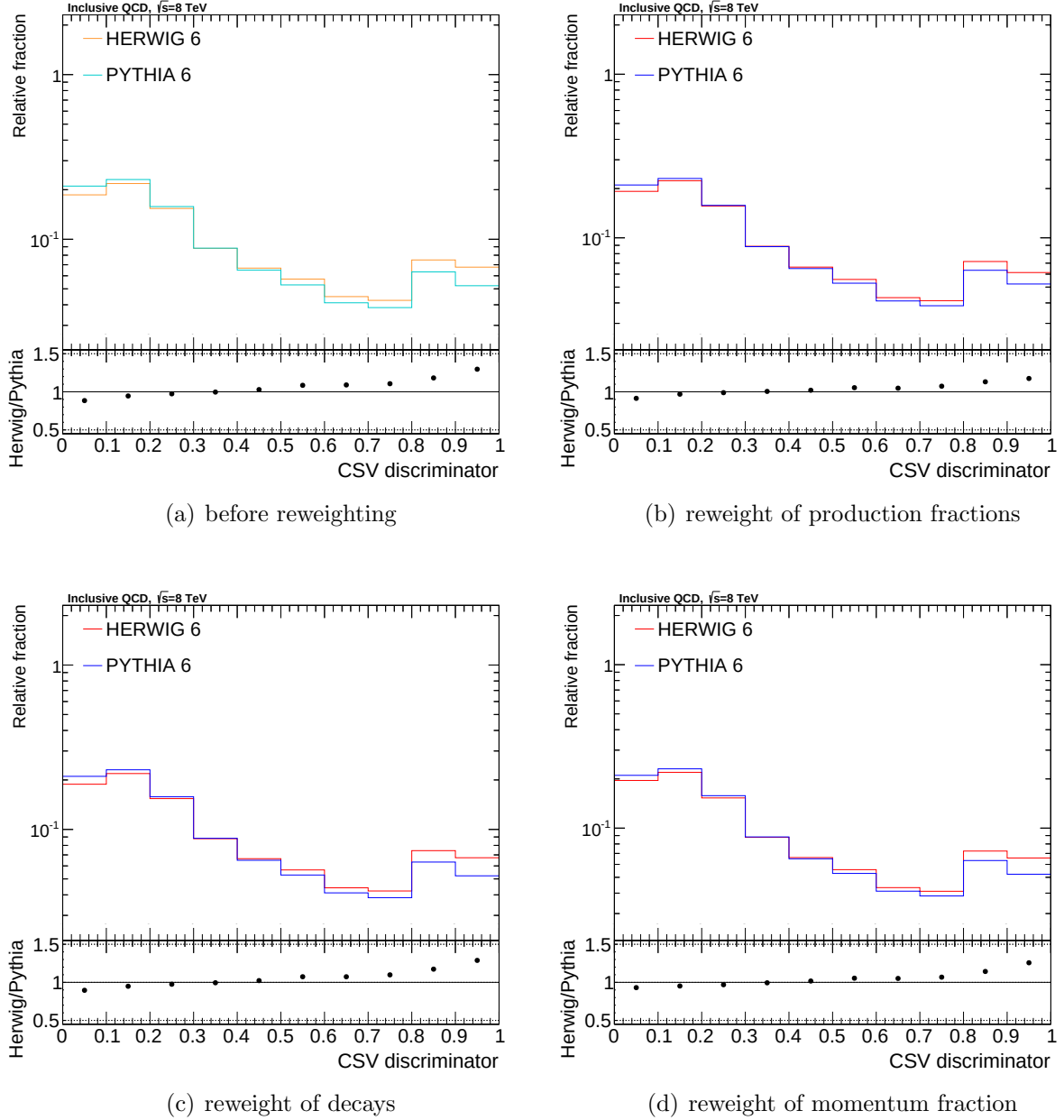


Figure 7.7.: CSV discriminator distribution of c-quark jets generated with PYTHIA 6 and c-quark jets generated with HERWIG 6 (a) before reweighting, after reweighting HERWIG 6 (b) charmed hadron production fractions, (c) charged-daughter multiplicity of charmed hadrons, or (d) charmed hadron momentum fractions to PYTHIA 6 predictions. Multijet events were generated for a centre-of-mass energy of $\sqrt{s} = 8$ TeV proton-proton collisions.

7.3.2. Combined reweighting

Reweighting just one quantity does not remove the differences in the CSV discriminator distributions of c-quark jets between PYTHIA 6 and HERWIG 6. It has to be checked if reweighting charmed hadron fractions, charged daughter multiplicity of charmed hadrons and the charmed hadron momentum fractions together removes these differences. Otherwise there might be other sources which contribute.

Since hadron decays and hadron momentum fraction depend on the charmed hadron type, this dependency has to be taken into account, when calculating the weights. First, weights $w(\text{type}, \text{decay})$ are calculated depending on the charmed hadron type and the charged-daughter multiplicity. They are calculated such that afterwards distributions of c-quark jets generated with HERWIG 6 match the predicted PYTHIA 6 distributions. In the next step, this weight is applied and afterwards the residual weights $w(p_T^D/p_T^J)$ depending on the momentum fraction are calculated. Each c-quark jet generated with HERWIG 6 is reweighted by applying the weight w :

$$w = w(\text{type}, \text{decay}) \times w(p_T^D/p_T^J). \quad (7.2)$$

Figure 7.8 shows the CSV discriminator distributions of c-quark jets generated with PYTHIA 6 and HERWIG 6, respectively. The Figure 7.8(a) shows the distributions before reweighting and Figure 7.8(b) shows the distributions after reweighting HERWIG 6 c-quark jets to PYTHIA 6 distributions. After reweighting, differences between both CSV distributions are smaller, but there are still some small differences especially for high CSV discriminator values.

Figure 7.9 shows the performance of the CSV algorithm. Differences in the performance become smaller after reweighting. There are still differences in the performance after reweighting for small b-quark jet efficiencies, for working points with a tight selection on the CSV discriminator value, originating from the differences in the CSV discriminator distribution.

In conclusion, there are differences in the charmed hadron modelling between PYTHIA 6 and HERWIG 6, which lead to a different performance of the CSV algorithm for c-quark jets generated with HERWIG 6 or PYTHIA 6. Differences in modelling charmed hadron production fractions, charmed hadron decays and charmed hadron momentum fractions can explain most of the discrepancies in performance, but there have to be some other differences in the modelling to explain the remaining differences.

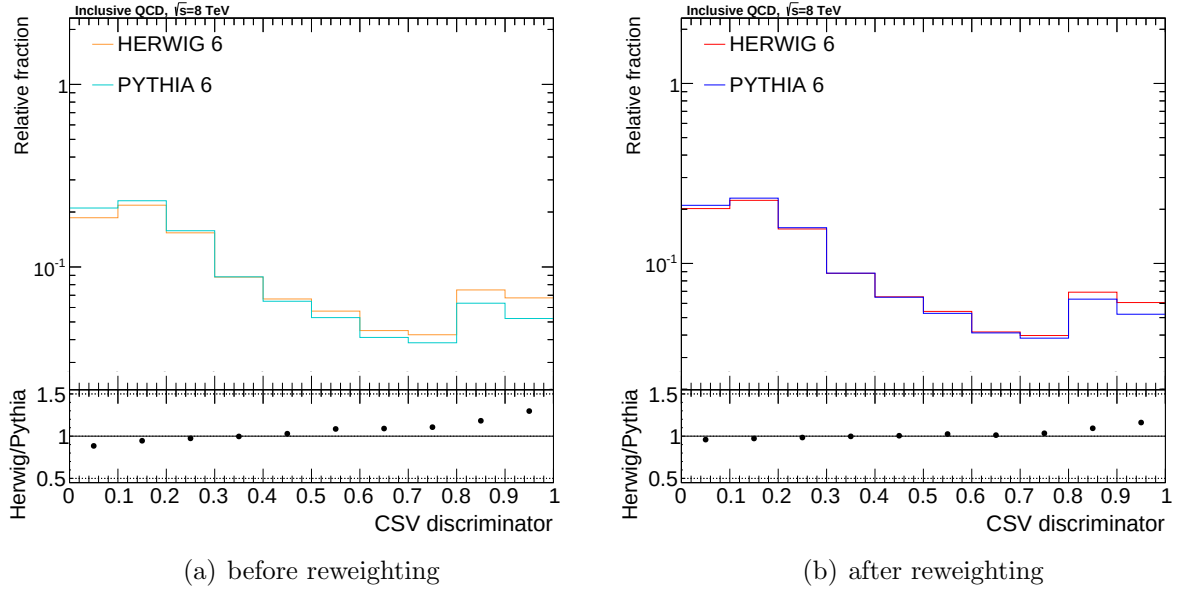


Figure 7.8.: CSV discriminator distribution of c-quark jets with a charmed hadron generated with PYTHIA 6 and HERWIG 6, respectively. Discriminator distributions of c-quark jets generated with HERWIG 6 are shown (a) before and (b) after reweighting charmed hadron production fractions, charged daughter multiplicities of charmed hadrons, and charmed hadron momentum fractions of c-quark jets generated with HERWIG 6 to PYTHIA 6 predictions. Multijet events were generated for proton-proton collisions at $\sqrt{s} = 8$ TeV.

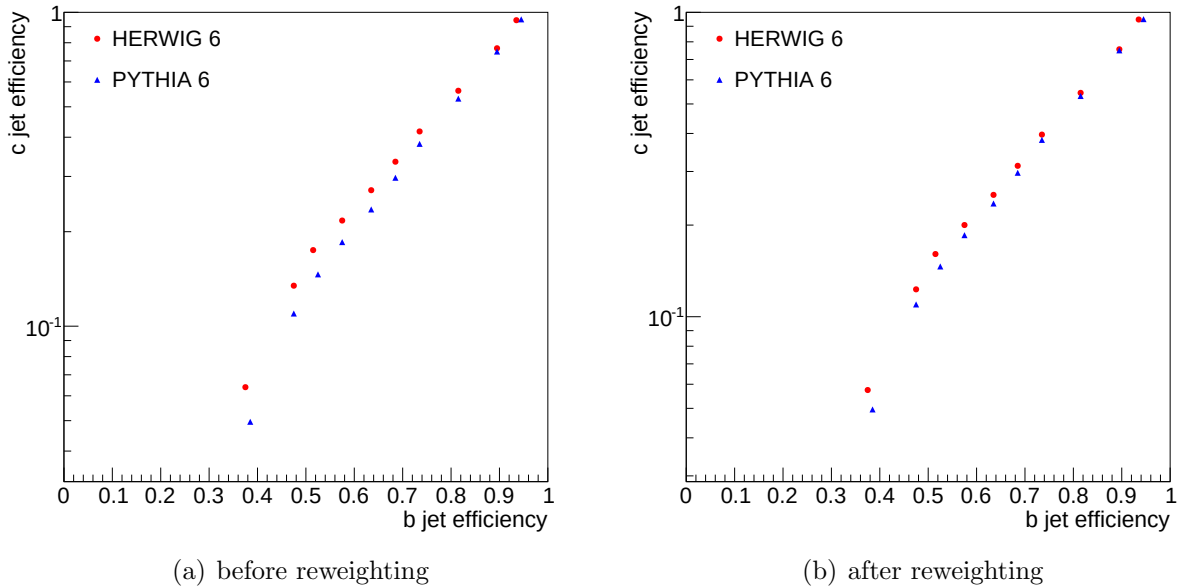


Figure 7.9.: Performance of the CSV algorithm for c-quark jets generated with PYTHIA 6 and HERWIG 6 (a) and (b) after reweighting HERWIG charmed hadron production fractions, charmed hadron decays and charmed hadron momentum fractions to PYTHIA distributions. QCD multijet events were generated for proton-proton collisions at $\sqrt{s} = 8$ TeV.

7.4. Reweighting to the measured values

In the previous section, it has been discussed that differences in heavy flavour modelling between PYTHIA 6 and HERWIG 6 affect the performance of the CSV algorithm for c -quark jets. Without measurements, it is difficult to judge, which model can describe observed distributions better. Thus, it is studied in this section which heavy flavour modelling can better describe observed distributions and values.

Charmed hadron production fractions, charged-daughter multiplicities, and charmed hadron momentum fractions are compared to observed values. There had not been explicit measurements of the discussed observables at the LHC, therefore observables distributions from other experiments with a lower centre-of-mass energy were used for this study. Thus, it is studied which of both MC generators describes charmed-hadron measurements best, but not which of them can describe charmed-hadrons in proton-proton collisions at $\sqrt{s} = 8$ TeV best. Distributions are reweighted to the measured ones and it is checked if there are still differences in the performance of the CSV algorithm for c -quark jets between PYTHIA 6 and HERWIG 6. If there are no differences left after reweighting to measurements, updating the settings of both MC generators or reweighting these quantities would resolve these discrepancies. Otherwise, a systematic uncertainty for differences in charmed-hadron modelling has to be taken into account when estimating the scale factor uncertainties of c -quark jets.

7.4.1. Charmed hadron production fractions

Charmed hadron production fractions were measured by the ZEUS Collaboration at HERA [122]. HERA collided electrons or positrons with an energy of $E_e = 27.5$ GeV and protons with an energy of $E_p = 920$ GeV leading to a centre-of-mass energy of:

$$E_{\text{CMF}} = \sqrt{s} = \sqrt{4E_e E_p} \cong 318 \text{ GeV}. \quad (7.3)$$

HERA had a smaller centre-of-mass energy than the LHC, but the other measurements of charmed hadrons were done with lepton colliders or at smaller centre-of-mass energies. The HERA measurements are closest to the conditions at the LHC, therefore the ZEUS measurements are used.

Table 7.5 lists the charmed hadron production fractions measured with the ZEUS detector, or predicted by PYTHIA 6 and HERWIG 6, respectively. Charmed hadron production fractions measured by the ZEUS Collaboration are between the values predicted by PYTHIA 6 and HERWIG 6, therefore the HERA charmed hadron production fraction can be compared to the LHC production fractions. Figure 7.10 shows the values visually.

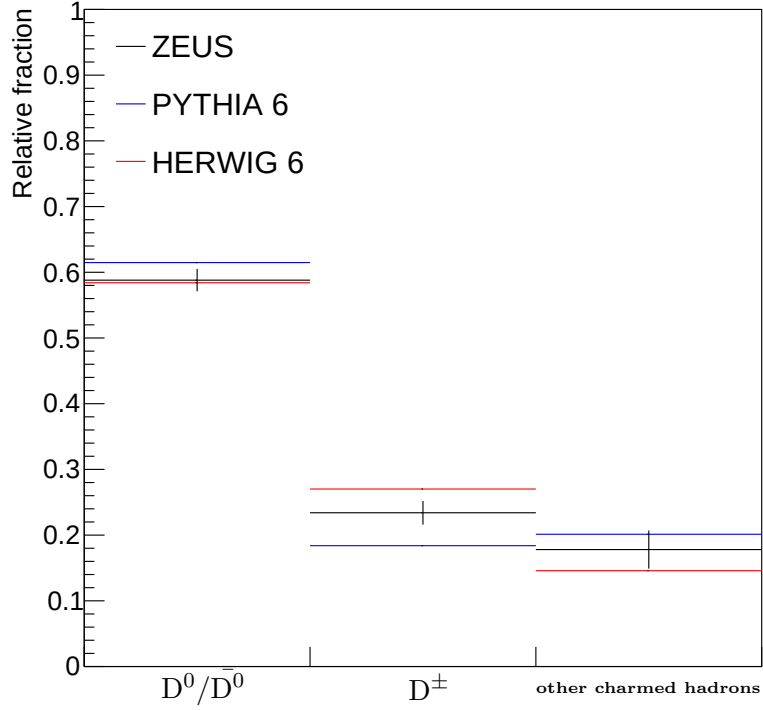


Figure 7.10.: Charmed hadron production fractions measured with the ZEUS detector [122] at $\sqrt{s} = 300$ GeV compared to production fractions predicted by PYTHIA 6 and HERWIG 6 for $\sqrt{s} = 8$ TeV.

Both predictions differ from the measured charmed hadron production fractions with all measured ZEUS values are between by PYTHIA 6 and HERWIG 6 predicted production fractions. For this reason, it is not possible to judge which modelling describes the observed values better.

Figure 7.11 shows the CSV discriminator distribution for c -quark jets with a charmed hadron predicted by PYTHIA 6 and HERWIG 6 before and after reweighting to ZEUS values. This has a negligible effect on the CSV discriminator distribution as it can be seen in 7.11(a) and (b), but it reduces the differences between CSV discriminator distribution predicted by PYTHIA 6 and HERWIG 6 as shown in Figure 7.11(c) and (d).

Table 7.5.: Charmed hadron production fractions measured with the ZEUS detector [122] at $\sqrt{s} = 300$ GeV compared to production fractions predicted by PYTHIA 6 and HERWIG 6 for $\sqrt{s} = 8$ TeV.

Type	ZEUS	PYTHIA	HERWIG
$D^\pm$	$(23.4^{+1.8}_{-1.7})\%$	$(18.39 \pm 0.06)\%$	$(27.01 \pm 0.05)\%$
$D^0/\bar{D}^0$	$(58.8^{+1.7}_{-1.0})\%$	$(61.48 \pm 0.12)\%$	$(58.40 \pm 0.08)\%$
Other charmed hadrons	$(17.8^{+2.5}_{-2.9})\%$	$(20.13 \pm 0.06)\%$	$(14.58 \pm 0.03)\%$

7. Impact of heavy flavour modelling on the misidentification of c -quark jets

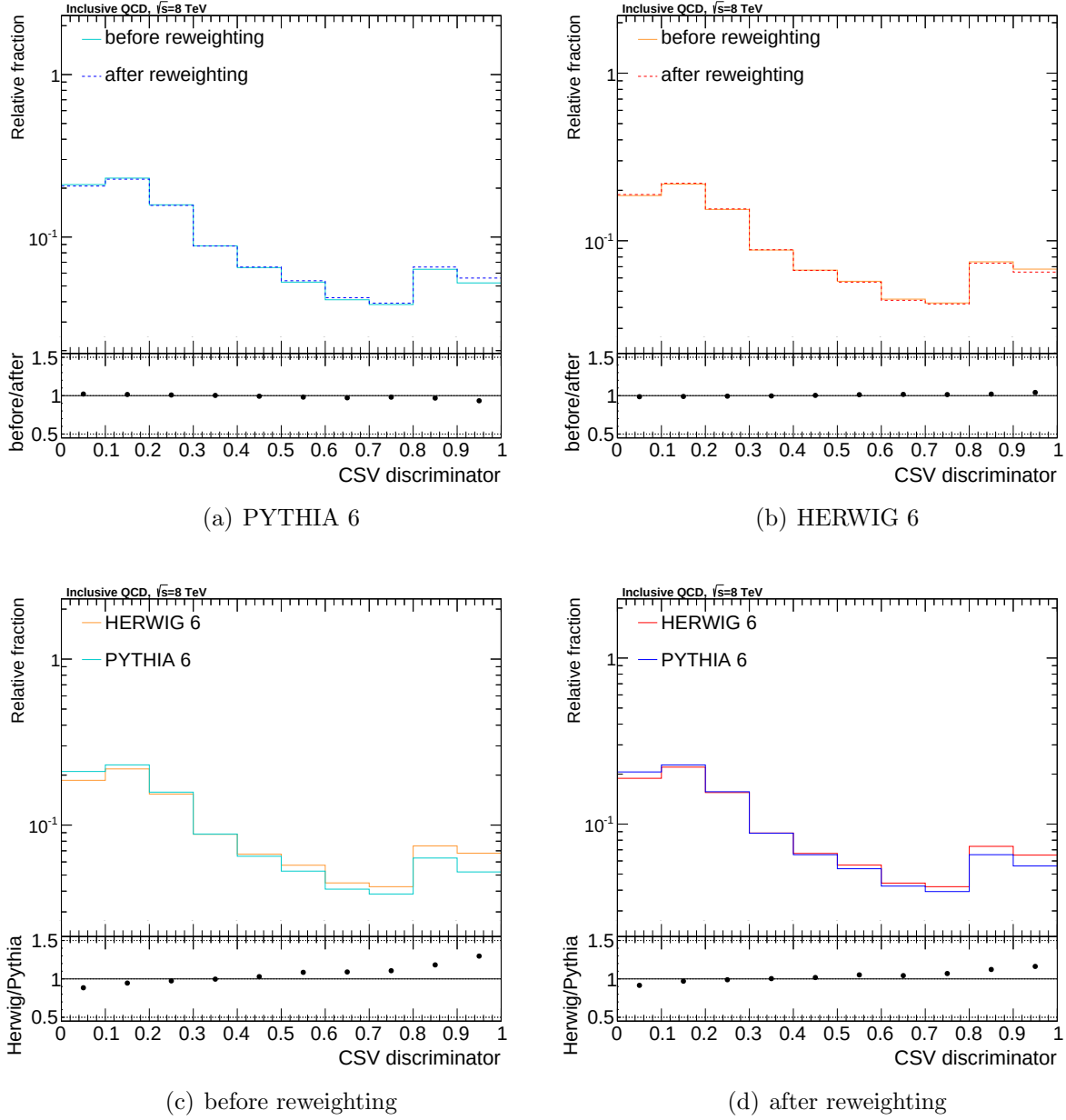


Figure 7.11.: CSV discriminator distribution of c -quark jets before and after reweighting charmed hadron production fractions predicted by (a) PYTHIA 6 and (b) HERWIG 6 at $\sqrt{s} = 8$ TeV to ZEUS values. Predicted CSV discriminator distributions are compared (c) before and (d) after reweighting.

Table 7.6.: Measured charged-daughter multiplicity of $D^0/\bar{D}^0$ mesons taken from the particle summary table of the particle data group [120]. Directly measured values are labelled PDG and the other probability values are calculated using exclusive decays.

	0	2	4	6
PDG	$(15 \pm 6\%)$	$(70 \pm 6)\%$	$(14.5 \pm 0.5)\%$	$(6.4 \pm 1.3) \times 10^{-4}$
Exclusive decays	9.4%	71.6%	18.8%	0.1%

7.4.2. Charmed hadron decays

Charged-daughter multiplicities are used to reweight the charmed hadron decays, since the misidentification probability mainly depends on the track multiplicity and on involved decay products. Furthermore, counting just the number of charged decay products is easier than considering each decay. In addition, particles from charmed-hadron decays do not have to be identified.

Charmed hadron charged-daughter multiplicities had not been measured at the LHC, when this study was done, therefore other measurements were considered. The particle data group [120] gives these values directly only for $D^0/\bar{D}^0$ mesons in its particle summary table. Charged-daughter multiplicities for remaining charmed hadrons can be calculated from their particle summary tables, which are giving an overview of available experimental measurements. These measurements have been done at experiments with a lower centre-of-mass energy than the one at the LHC. As in paragraph 5.2 discussed, particle summary tables can be incomplete or inconsistent within experimental precision, and all branching fractions do not have to be adding up to unity. Thus, choices have to be made when adapting summary tables and double counting has to be avoided.

This method is validated by using directly measured values for $D^0/\bar{D}^0$ mesons and comparing them to calculated values. Most important decays are considered and they are classified according to their charged-daughter multiplicity. Then, all branching fractions are added up for each multiplicity. The used branching fractions add up to about $\sim 75\%$. Finally, probability values are normalised, so that they add up to unity. Table 7.6 shows direct probability values for $D^0/\bar{D}^0$ mesons and the result from using exclusive decays of the particle summary table. About 70% of $D^0/\bar{D}^0$ mesons decay into two charged particles and neutral particles. The use of the value from exclusive decays gives the same value as the direct measurement within experimental precision. Also values of the other two next important decay modes with zero and four charged daughters are comparable. There are differences for the last decay mode with six charged particles, but this decay mode is negligible. In conclusion, using exclusive decays for extracting charged-daughter multiplicities is a good method, when there are no direct measurements.

7. Impact of heavy flavour modelling on the misidentification of c-quark jets

Charged hadron multiplicity can be calculated for $D^\pm$ mesons using exclusive decays taken from the particle summary table of the particle data group [120]. However, this method is not practicable for the *other charmed hadrons*, since several charmed hadrons are combined to this category, neglecting less important hadrons. Combining exclusive decays of one particle table is quite complicated, combining several of these tables, taking also production fraction into account is a lot of work and imprecise. This category is dominated by charmed strange mesons D_s^+/D_s^- , which have a smaller mass than charmed baryons as it can be seen in Table 7.1. Also the next important charmed strange baryons, Ξ_c^+/Ξ_c^- , are charged. About 18% of all c-quark jets contain a hadron of this category as shown in Table 7.5. As charge is conserved, there have to be mainly odd numbers of charged daughters. The cases where a hadron of this category has no charge, and therefore decays into an even number of charged particles, should be negligible in comparison to all other c-quark jets containing $D^0/\bar{D}^0$ mesons, $D^\pm$ mesons, or charged *other charmed hadrons*. Figure 7.2 shows the simulated charged-daughter multiplicity distribution of *other charmed hadrons*. It can be seen that most of these hadrons decay into one or three charged particles and neutral particles. Comparing these values to the charged-daughter multiplicity distribution of $D^\pm$ mesons, it can be noticed that they are comparable, therefore it should be a good approximation to use this values instead of calculating them from particle tables of *other charmed hadrons*.

Table 7.7.: Measured charmed hadron charged-daughter multiplicity values used for reweighting. Values of $D^0/\bar{D}^0$ mesons have been taken directly from [120]. Values of both other categories have been calculated using exclusive decays in the particle summary table of $D^\pm$ mesons [120].

Type	0	1	2	3	4	5	6
$D^0/\bar{D}^0$	$(15 \pm 6)\%$	—	$(70 \pm 6)\%$	—	$(14.5 \pm 0.5)\%$	—	$(6.4 \pm 1.3) \times 10^{-4}$
$D^\pm$	—	55.7%	—	43.3%	—	1.1%	—
Other charmed hadrons	—	55.7%	—	43.3%	—	1.1%	—

Charged-daughter multiplicities used for doing the reweighting of charmed hadron decays are listed in Table 7.7. Figure 7.12 compares these values with PYTHIA 6 and HERWIG 6 for each charmed hadron type. Both MC generator predictions do not fit to particle data group values, but PYTHIA 6 charmed hadron decay modelling seems to be closer to measurements.

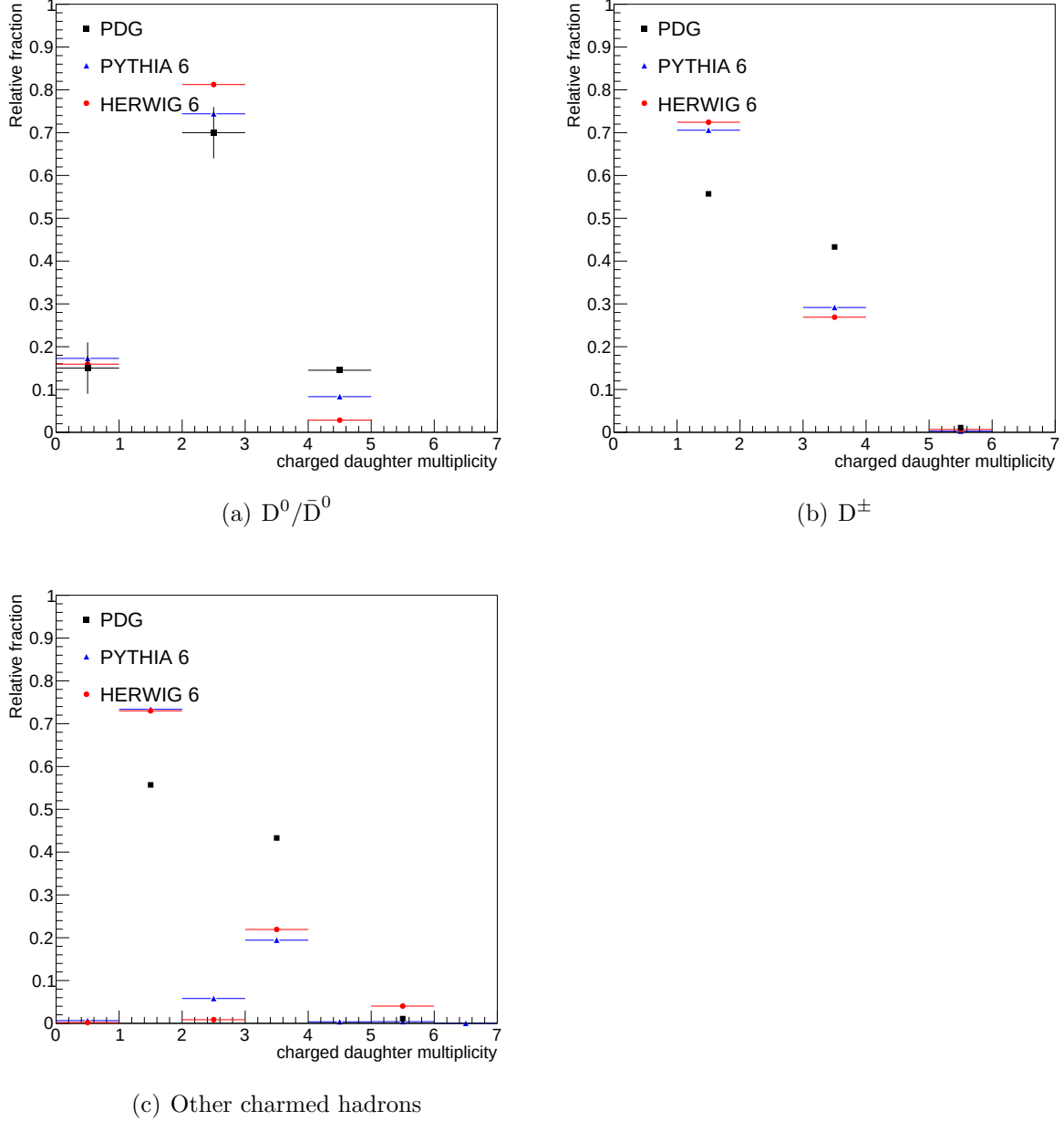


Figure 7.12.: Comparison of charmed hadron charged-daughter multiplicities measured by particle data group, or simulated by PYTHIA 6 or HERWIG 6, respectively. Multijet events were generated for a centre-of-mass energy of $\sqrt{s} = 8$ TeV. Charged-daughter multiplicities are shown for each charmed hadron type: (a) $D^0/\bar{D}^0$ mesons, (b) $D^\pm$ mesons, and (c) *other charmed hadrons*.

7. *Impact of heavy flavour modelling on the misidentification of c-quark jets*

Figure 7.13 shows CSV discriminator distributions before and after reweighting charmed hadron decays to measured values. Differences between before and after reweighting can be seen in (a) for c-quark jets generated with PYTHIA 6 and in (b) for c-quark jets generated with HERWIG 6. Mainly high discriminator values are affected. It can be seen in (c) and (d) that differences between PYTHIA 6 and HERWIG 6 in the CSV discriminator distribution are not affected by reweighting.

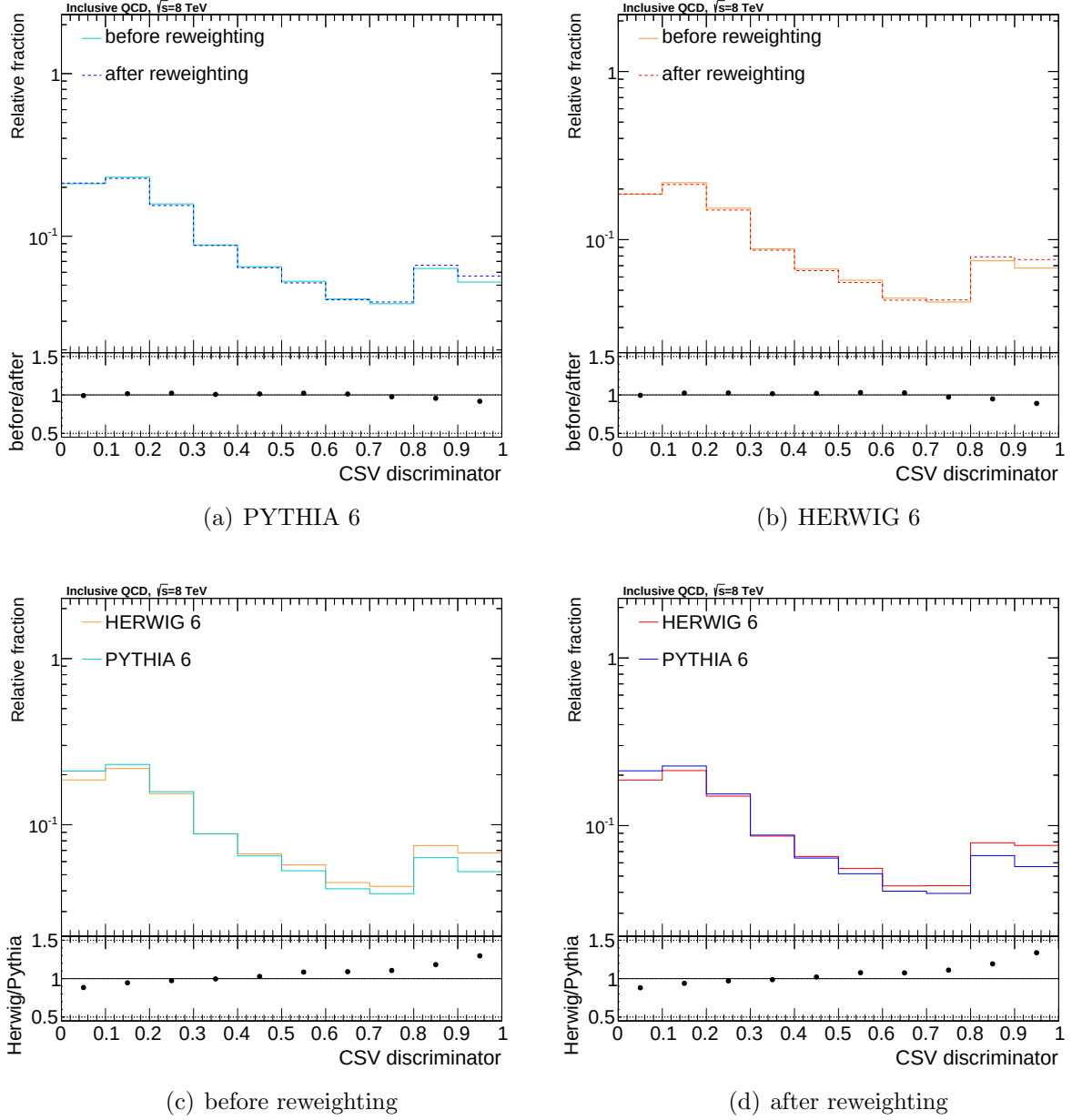


Figure 7.13.: CSV discriminator distributions of c-quark jets generated with PYTHIA 6 and HERWIG 6. Multijet events are generated for a centre-of-mass energy of $\sqrt{s} = 8$ TeV. Simulated charmed hadron decays are reweighted to measured charged daughter multiplicities. CSV distributions are shown (a) for PYTHIA before and after reweighting, (b) for HERWIG before and after reweighting, (c) for PYTHIA and HERWIG before reweighting and (d) for PYTHIA and HERWIG after reweighting.

7.4.3. Charmed hadron momentum fraction

In this section, charmed hadron momentum fractions p_T^D/p_T^J are investigated. In section 5, fragmentation functions were discussed, which give the probability distributions of the momentum fraction z carried by the hadron from the corresponding quark momentum. They have several free parameters which can be adjusted by fitting the measurements. One of these fragmentation functions with just one free parameter is the Peterson/SLAC fragmentation function [64]. This function is defined for heavy flavour as:

$$f(z) \propto \frac{1}{z \left(1 - \frac{1}{z} - \frac{\epsilon_c}{1-z}\right)^2}, \quad (7.4)$$

which depends on the free parameter ϵ_c , which can be described as function of the position of the maximum of this fragmentation function $\langle z \rangle$:

$$\epsilon_c = \frac{(\langle z \rangle - 1)^2}{\langle z \rangle}. \quad (7.5)$$

The free parameter depends on the hadron type. Bottom hadrons have the maximum of the momentum fraction distribution at high values and charmed hadrons have their maxima at smaller momentum fractions than bottom hadrons. It also depends on the charmed hadron type. Figure 7.14 shows the Peterson/SLAC fragmentation function for different positions of the maximum. Its distribution has a smaller width for maxima at larger z values.

Jet and quark momenta should be comparable, therefore charmed hadron momentum fraction distributions should have a similar distribution as the Peterson/SLAC fragmentation function. Position of the maximum and ϵ_c have been measured by the CLEO Collaboration in e^+e^- collisions at a centre-of-mass energy of 10.5 GeV [123].

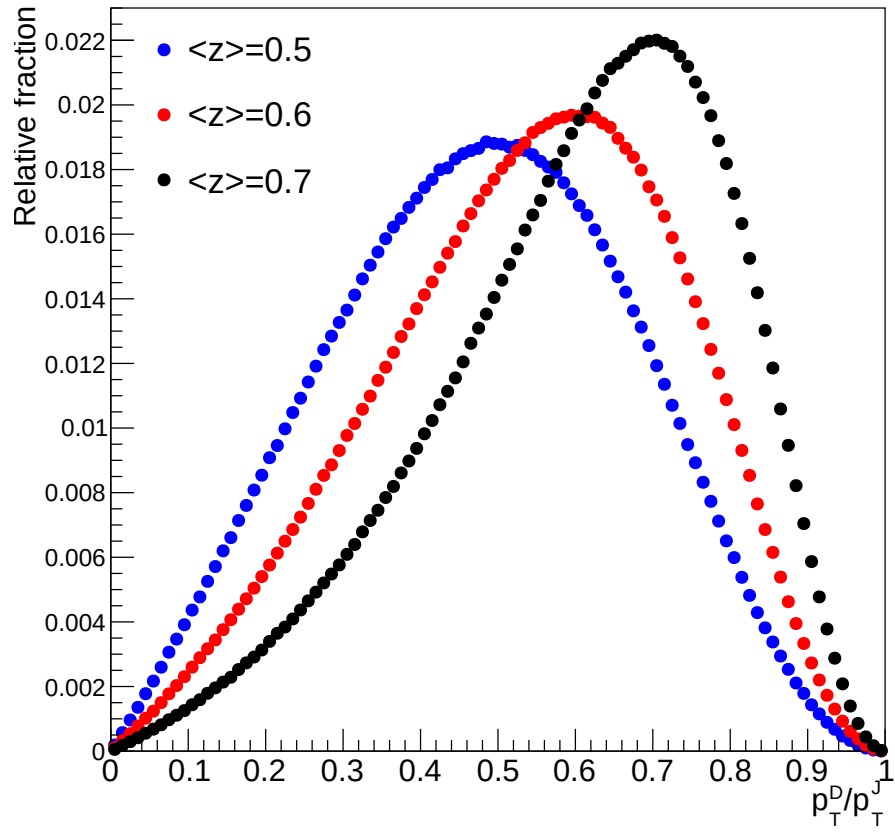


Figure 7.14.: The Peterson/SLAC fragmentation function for heavy flavour is shown for different $\langle z \rangle$ values: 0.5, 0.6, and 0.7.

7. Impact of heavy flavour modelling on the misidentification of c -quark jets

Table 7.8 lists the maxima positions of the Peterson fragmentation function for each charmed hadron type. Measured values by CLEO and simulated values for PYTHIA 6 and HERWIG 6 are listed in the table. A fit of the Peterson/SLAC fragmentation function to the generated charmed hadron momentum fraction distributions was done for each MC generator, where $\langle z \rangle$ is the maximum of the best fit. There are no measured values for the *other charmed hadrons*. Differences between charmed hadron types are quite small. Both MC generators use similar values for $D^\pm$ mesons and *other charmed hadrons*. HERWIG 6 uses a harder fragmentation than PYTHIA 6 and the CLEO results, therefore PYTHIA can better describe the measurements.

Table 7.8.: Maximum position of the Peterson fragmentation function for charmed hadron momentum fractions, for CLEO, PYTHIA 6 and HERWIG 6. The maximum of the momentum fraction depends on the charmed hadron type. Position of the maximum and ϵ_c have been measured by the CLEO Collaboration in e^+e^- collisions at a centre-of-mass energy of 10.5 GeV [123]. Multijets were generated for proton-proton collisions at $\sqrt{s} = 8$ TeV.

Type	CLEO $\langle z \rangle$	PYTHIA fitted $\langle z \rangle$	HERWIG fitted $\langle z \rangle$
$D^0/\bar{D}^0$	0.570	0.56	0.62
$D^\pm$	0.582	0.53	0.63
Other charmed hadrons	—	0.54	0.63

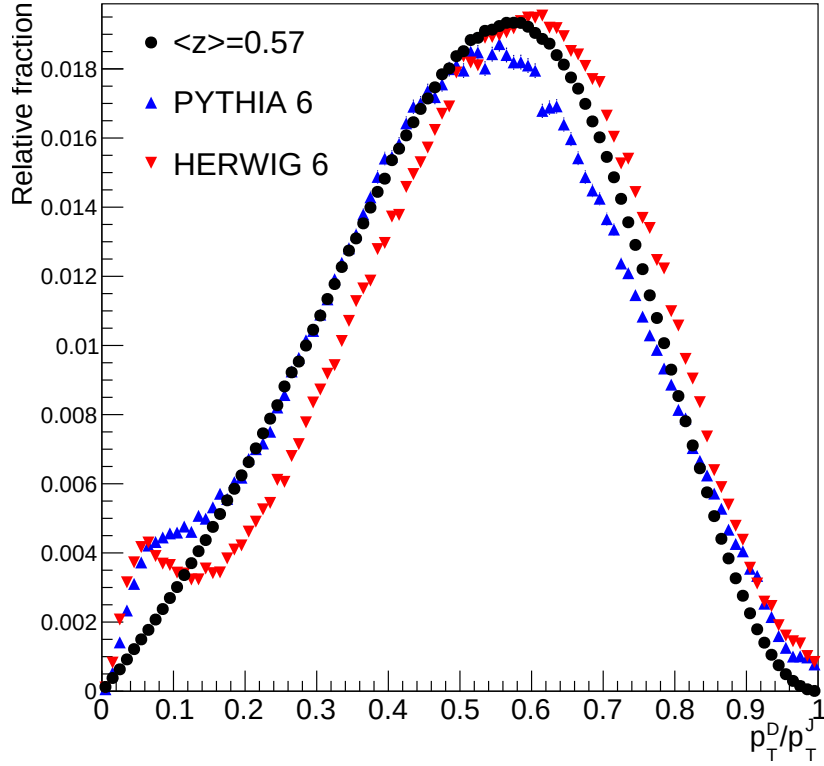


Figure 7.15.: Momentum fraction distribution of $D^0/\bar{D}^0$ mesons. The Peterson fragmentation function with maximum at $\langle z \rangle = 0.57$ measured by the CLEO Collaboration in e^+e^- collisions at $\sqrt{s} = 10.5$ GeV [123]. Multijet events were generated for proton-proton collisions at $\sqrt{s} = 8$ TeV using PYTHIA 6 and HERWIG 6, respectively.

Figure 7.15 shows the Peterson fragmentation function using the free parameter for D^0 mesons measured by the CLEO Collaboration in comparison to the momentum fraction distribution of c-quark jets containing a $D^0/\bar{D}^0$ meson generated by PYTHIA 6 and HERWIG 6, respectively. The Peterson fragmentation function describes both distributions well, except for small momentum fractions, where a small peak for both generators is visible. This extra events could come from mismatching. Charmed hadrons and c-quark jets are matched by ΔR . A requirement on the momentum of the matched charmed hadron to the jet momentum could reduce the amount of mismatched events. But just about 1% of the events have a wrong match, therefore this should have no effect on the CSV discriminator shape. These events have lower CSV discriminator values than the loose CSV working point, thus they are not responsible for the differences in the performance of the CSV algorithm between PYTHIA 6 and HERWIG 6, since these differences are quite small for low CSV discriminator values. In conclusion, it should be fine to reweight generated momentum fraction distributions to Peterson fragmentation functions using parameters measured by the CLEO Collaboration.

7. Impact of heavy flavour modelling on the misidentification of c -quark jets

Both MC generators are reweighted to the CLEO values. It is assumed that the momentum fraction distribution of the *other charmed hadrons* has a maximum at 0.58 as for $D^\pm$ mesons, since this value has not been measured by the CLEO Collaboration and both MC generators use similar values for $D^\pm$ mesons and *other charmed hadrons*. Figure 7.16 shows the CSV distributions of c -quark jets generated with PYTHIA 6 and HERWIG 6, respectively. Distributions are shown before and after reweighting charmed hadron momentum fractions. It can be seen in Figure 7.16(a) and (b) that reweighting charmed hadron momentum fractions has no effect on the shape of the CSV distribution, therefore reweighting charmed hadron momentum fractions to CLEO values does not reduce differences between the MC generators in the CSV discriminator distribution.

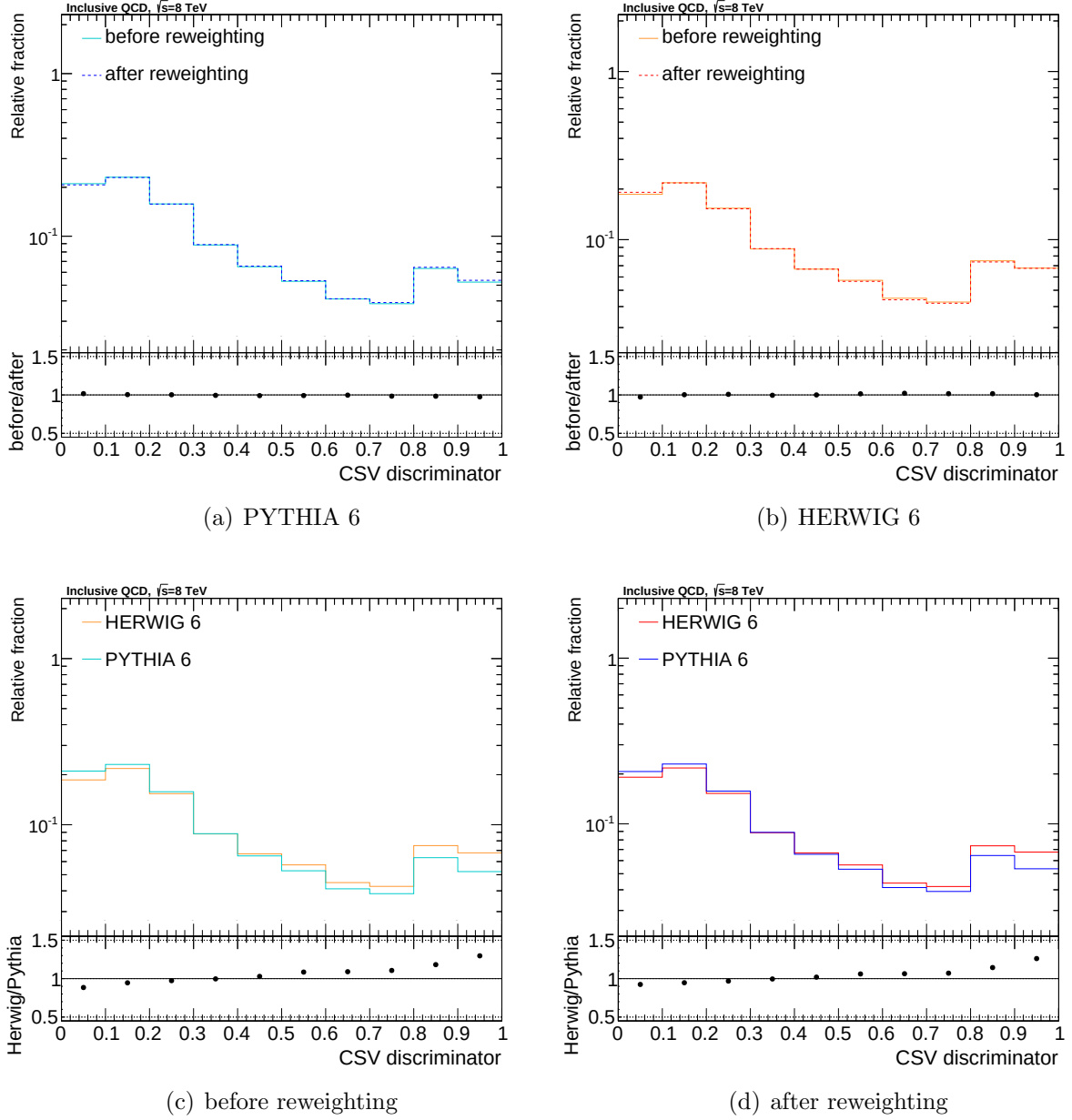


Figure 7.16.: CSV discriminator distributions of c-quark jets generated with PYTHIA 6 and HERWIG 6, respectively. Multijets were generated for a centre-of-mass energy of $\sqrt{s} = 8$ TeV. Charmed hadron momentum fractions are reweighted to CLEO values. The CSV distribution is shown (a) for PYTHIA 6 before and after reweighting, (b) for HERWIG 6 before and after reweighting, (c) for PYTHIA 6 and HERWIG 6 before reweighting, and (d) for PYTHIA 6 and HERWIG 6 after reweighting.

7.4.4. Remaining differences between HERWIG and PYTHIA

Generated charmed hadron production fractions, charged-daughter multiplicities and momentum fractions were compared to measured distributions. Measurements by the ZEUS Collaboration were used for charmed hadron production fractions. PDG values were used for charmed hadron decays using the charged-daughter multiplicity. Charmed hadron momentum fractions can be described by the Peterson/SLAC fragmentation function for heavy hadrons using measurements by the CLEO Collaboration. Neither PYTHIA nor HERWIG describe the measurements.

A reweighting was done in each observable for c-quark jets generated with PYTHIA 6 and HERWIG 6, respectively. Reweighting to measured charmed-hadron production-fractions values has the largest impact on the CSV discriminator shape. Reweighting charmed hadron decays has an effect on the high CSV discriminator values, whereas reweighting charmed momentum fractions has no effect on the discriminator shape.

Reweighting c-quark jets generated with HERWIG 6 to PYTHIA 6 values in these observables does not remove all differences between both MC generators. It has to be checked if reweighting c-quark jets generated by PYTHIA 6 or HERWIG 6 to measured values in all three observables removes differences between both MC generators in the CSV discriminator shape. There could be still remaining differences which would come from other sources. Figure 7.17 compares the CSV discriminator distribution for c-quark jets generated with PYTHIA 6 and c-quark jets generated with HERWIG 6. Distributions are shown (a) before and (b) after simultaneous reweighting charmed hadron production fractions to ZEUS values, charged daughter decays to PDG values and momentum fractions to CLEO values. Most differences are reduced, especially in the medium region, but there are still differences left for high and small discriminator values.

In Figure 7.18, CSV discriminator distributions for c-quark jets generated with PYTHIA 6 and HERWIG 6, respectively, are shown. First, CSV discriminator distributions are compared after reweighting HERWIG 6 to PYTHIA 6 charmed hadrons distributions as shown in Figure 7.18(a). Then, these distributions are shown for reweighting charmed hadrons to measured values in Figure 7.18(b).

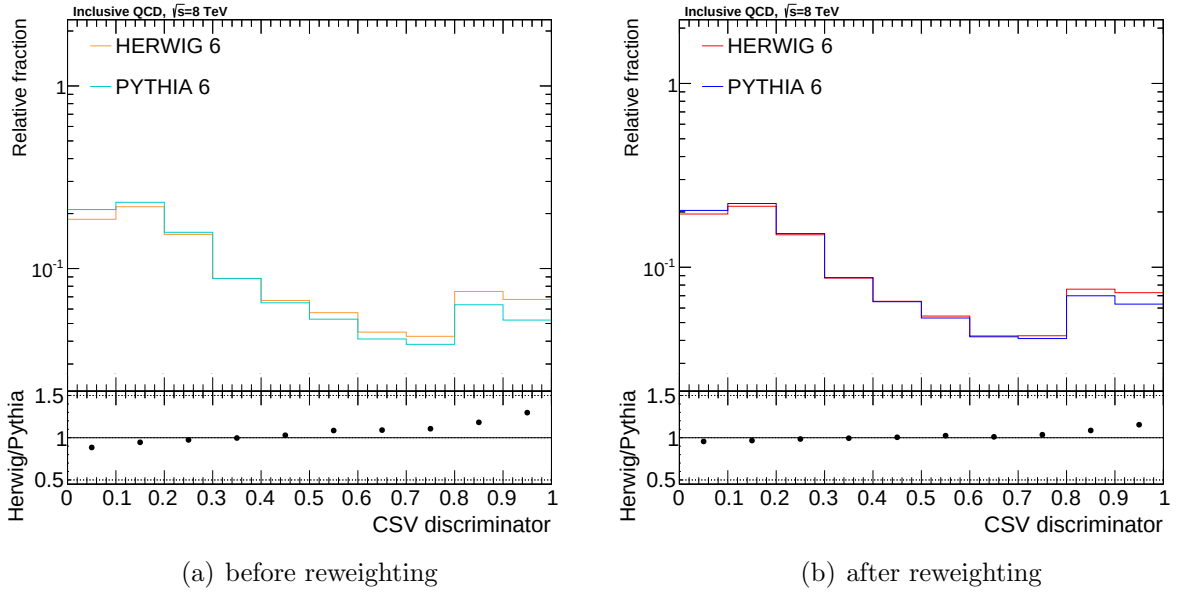


Figure 7.17.: CSV discriminator distribution after reweighting c-quark jets generated with PYTHIA 6 or HERWIG 6 to charmed hadron production fractions measured by the ZEUS Collaboration, to PDG charged-daughter multiplicities, and momentum fractions measured by the CLEO Collaboration. Multijets were generated for proton-proton collisions at $\sqrt{s} = 8$ TeV. Distributions are shown (a) before and (b) after reweighting.

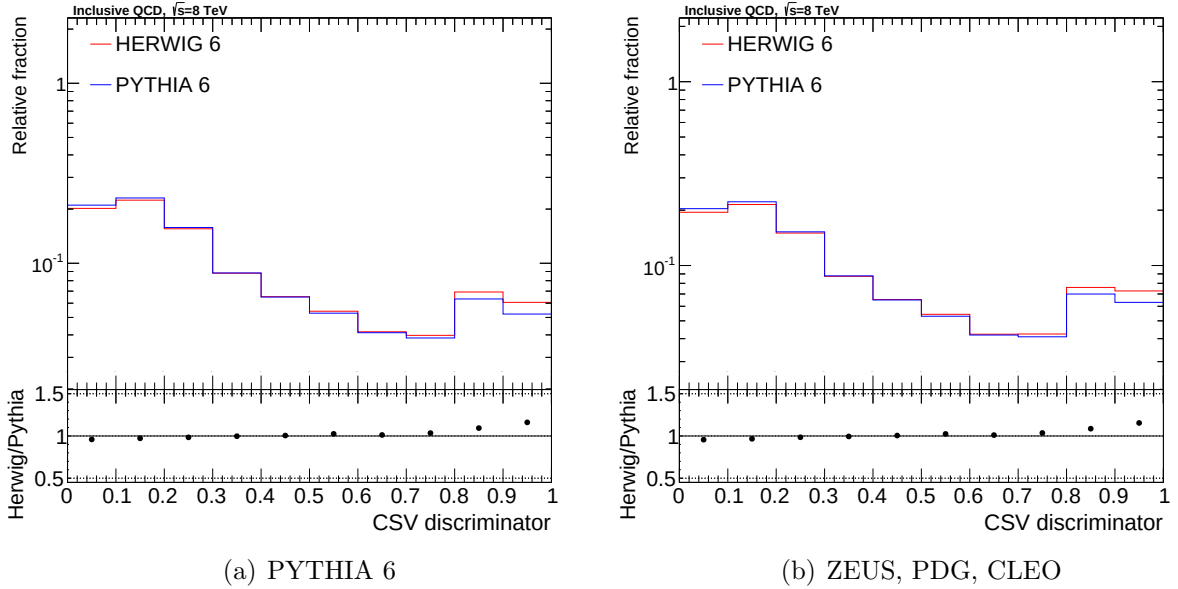


Figure 7.18.: CSV discriminator distribution of c-quark jets generated with PYTHIA 6 and HERWIG 6, respectively. Multijets were generated for proton-proton collisions at $\sqrt{s} = 8$ TeV. Distributions are shown after reweighting charmed hadron production fractions, charged-daughter multiplicities and momentum fractions (a) of c-quark jets generated with HERWIG 6 to PYTHIA 6 values (b) of c-quark jets generated with PYTHIA 6 and HERWIG 6 to ZEUS, PDG and CLEO values.

7. Impact of heavy flavour modelling on the misidentification of c-quark jets

First, it was shown in Section 7.3 that there are differences in modelling charmed hadron production fractions, decays and momentum fractions between PYTHIA 6 and HERWIG 6. Differences between the CSV discriminator distributions of c-quark jets between both generators are mainly coming from differences in modelling the charmed hadron production fractions. Differences in modelling the other two observables have also an effect on the CSV discriminator distribution. There are some unknown differences in modelling charmed hadrons, which have to be studied.

In this section, it was shown that there are differences between modelling charmed hadron production, charmed hadron decays and charmed momentum fractions of PYTHIA 6 and HERWIG 6 to measured values of ZEUS, PDG and CLEO. Neither PYTHIA 6 nor HERWIG 6 describes measured observables distributions. Reweighting charmed hadrons in these three observables shows that modelling charmed hadron production fractions has the largest effect on differences between the CSV discriminator distribution of c-quark jets generated with PYTHIA 6 and HERWIG 6, respectively. Differences in modelling the other two observables have also a small effect on the CSV discriminator distribution. There are also some unknown differences in modelling charmed hadrons, which have the same effect on differences in the CSV discriminator distribution for c-quark jets between both generators.

7.4.5. Effect on the misidentification probabilities of c-quark jets

It was shown that reweighting simulated charmed hadron distributions to measured distributions has an effect on the CSV discriminator shape. Different working points are used in analyses to select or remove jets identified as b-quark jets. Jets originating from c-quark jets are the main background. In this section it is presented how reweighting generated charmed hadron production fractions, charmed hadron decays and momentum fractions to measured values, changes the misidentification probability for each working point.

Table 7.9 lists the CSV discriminator for loose, medium and tight working points and the relative changes in the misidentification probability for c-quark jets generated with PYTHIA 6. Relative change is defined as difference in misidentification probability before and after reweighting divided by the misidentification probability before reweighting to measured values. This is done for each single reweighting and for doing all three reweightings together. It can be seen that doing the reweightings has the largest effect on the tight working point, especially for reweighting production fractions and decays. Remaining differences to HERWIG 6 after reweighting to ZEUS, PDG and CLEO values are also shown. These differences are of the same order as changes in misidentification probability before and after doing the reweighting.

Table 7.9.: Relative changes in misidentification probability for c-quark jets defined as difference in misidentification probability before and after reweighting divided by the probability before reweighting. Remaining differences to HERWIG 6 are similar defined after reweighting both MC generators to ZEUS, PDG and CLEO values. Multijet events were generated using PYTHIA 6 for proton-proton collisions at $\sqrt{s} = 8$ TeV.

Working point	Loose	Medium	Tight
CSV discriminator	> 0.244	> 0.679	> 0.898
Production fractions (ZEUS)	+1.5%	+4.0%	+6.9%
Charged-daughter multiplicity (PDG)	+1.3%	+5.7%	+9.4%
Fragmentation (CLEO)	+1.1%	+2.0%	+2.6%
ZEUS+PDG+CLEO	+4.4%	+13.0%	+21.0%
Remaining MC differences to HERWIG 6	-4.5%	-10.9%	-17.1%

7. Impact of heavy flavour modelling on the misidentification of c-quark jets

Table 7.10 lists results for c-quark jets generated with HERWIG 6. Reweighting charmed hadron decays has the largest effect on the misidentification probability. Doing the reweighting in all three observables has the largest effect on the tight working point. There are still differences to misidentification probabilities of c-quark jets generated with PYTHIA 6.

Table 7.10.: Relative changes in misidentification probability for c-quark jets defined as difference in misidentification probability before and after reweighting divided by the probability before reweighting. Remaining differences to HERWIG 6 are similar defined after reweighting both MC generators to ZEUS, PDG and CLEO values. Multijet events were generated using PYTHIA 6 for proton-proton collisions at $\sqrt{s} = 8$ TeV.

	Loose	Medium	Tight
CSV discriminator	> 0.244	> 0.679	> 0.898
Production fractions (ZEUS)	-1.0%	-2.5%	-3.9%
Charged-daughter multiplicity (PDG)	$+1.8\%$	$+7.5\%$	$+12.9\%$
Fragmentation (CLEO)	-0.3%	-0.8%	$+0.1\%$
ZEUS+PDG+CLEO	$+0.4\%$	$+3.7\%$	$+8.3\%$
Remaining MC differences to PYTHIA 6	$+4.3\%$	$+9.8\%$	$+14.6\%$

As has been shown in this section, reweighting charmed hadron production fractions, decays and momentum fractions to measured values does not remove differences between misidentification probabilities of c-quark jets generated with PYTHIA 6 or HERWIG 6. Thus, the performance of the CSV algorithm depends on whether jets were generated using HERWIG 6 or PYTHIA 6, therefore it has to be studied if these differences are covered by systematic uncertainties.

7.5. Systematic uncertainties on misidentification probability of c-quark jets

Since c-quark jets are the main background for the identification of b-quark jets and there are differences in the heavy flavour modelling of charmed hadrons which can affect the performance of the CSV algorithm, it has to be investigated if these differences are covered by the systematic uncertainties. This is especially important as a reweighting cannot remove all differences between PYTHIA 6 and HERWIG 6.

Figure 7.19 shows for each CSV working point the misidentification probability of c-quark jets generated with PYTHIA 6 and HERWIG 6, respectively. Efficiencies for c-quark jets are shown before and after reweighting charmed hadron production fractions, decays and momentum fractions to ZEUS, PDG and CLEO values. It can be seen that c-quark jets generated with PYTHIA 6 are more affected by reweighting than c-quark jets generated using HERWIG 6. Misidentification probabilities of c-quark jets generated with PYTHIA 6 and HERWIG 6 are closer to each other after doing this reweighting. Using PYTHIA 6 for generating c-quark jets leads to a lower misidentification probability.

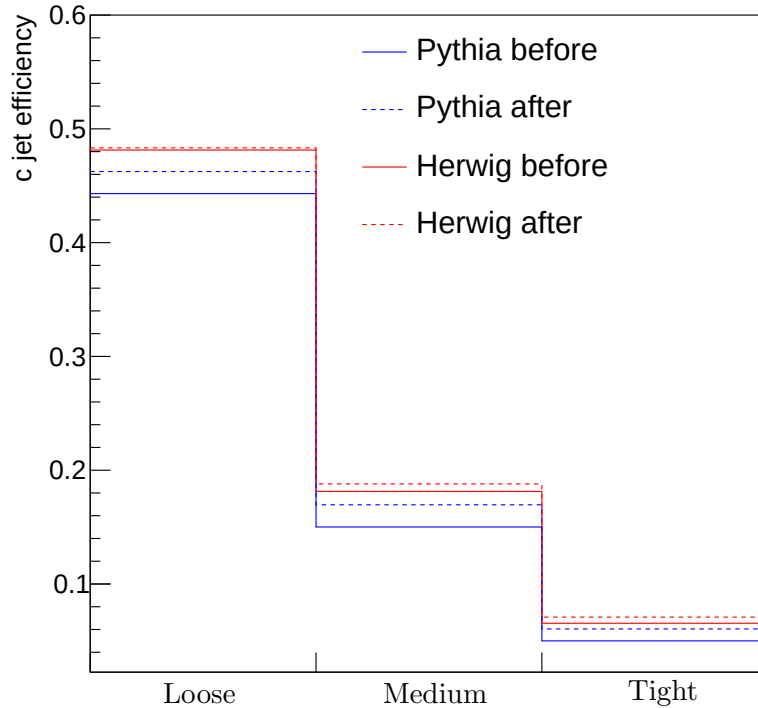


Figure 7.19.: Misidentification probability of c-quark jets for each CSV working point before (continuous line) and after (dashed line) reweighting to ZEUS, PDG and CLEO values. Multijets were generated with PYTHIA 6 (blue) and HERWIG 6 (red) for proton-proton collisions at $\sqrt{s} = 8$ TeV.

7. Impact of heavy flavour modelling on the misidentification of c-quark jets

Scale factors are applied on generated c-quark jets to correct for differences between measured CSV distributions and simulated ones. These scale factors were not measured for c-quark jets, but only for light flavour and b-quark jets. Measuring these scale factors for c-quark jets is more complicated, because it is difficult to get a validation sample with high purity. Since bottom hadrons and charmed hadrons have similar properties, b-quark jet scale factors are applied on c-quark jets. As the systematic uncertainties cannot be estimated precisely, systematic uncertainties twice as large as the b-quark jet scale factors are used instead.

It has to be studied if differences in heavy flavour modelling are covered by this systematic uncertainties. Especially, since there are differences between both MC generators which cannot be removed by correcting simulated charmed hadron observables to measured distributions. This is also important, as scale factors were derived for jets generated with PYTHIA 6 as reference. Since there are differences in the charmed hadrons, differences in the scale factors are expected.

Reweighting c-quark jets generated with PYTHIA 6 to measured values has a larger effect on the misidentification probability than reweighting c-quark jets generated with HERWIG 6. Since single reweightings for HERWIG 6 affect the discriminator in different directions as can be seen in Table 7.10. All single reweightings for PYTHIA 6 reweight the CSV discriminator shape to higher values as it can be seen in Table 7.9. After reweighting, c-quark jets generated with PYTHIA 6 have a higher misidentification probability. This means that the c-quark jet background could be underestimated.

For previous CMS publications, systematic uncertainties on scale factors for $\sqrt{s} = 8$ TeV were derived binned in jet p_T [124]. The total uncertainty on events of the multijet sample passing a working point can be calculated by using these scale factor uncertainties. Events are counted in each p_T bin and this number is multiplied by the corresponding scale factor uncertainty for this p_T bin. The total systematic uncertainty is given by the sum of the momentum dependent uncertainties. This uncertainty is normalised dividing by the total number of events. In the following, scale factor uncertainty is defined as this total uncertainty. Scale factor uncertainties for PYTHIA 6 were also used for calculating the HERWIG 6 scale factor uncertainty.

Table 7.11 lists the results for c-quark jets generated with PYTHIA 6. They are listed for each working point of the CSV algorithm. Relative changes when reweighting to measured values and remaining differences to c-quark jets generated by HERWIG 6 are listed. Also the relative difference to HERWIG 6 before reweighting and the scale factor uncertainty of this sample are listed. Table 7.12 shows a similar table for HERWIG 6.

Table 7.11.: Relative changes in misidentification probability for each CSV working point after reweighting c -quark jets generated by PYTHIA 6 to ZEUS, PDG and CLEO values. Remaining differences (after reweighting) and relative differences (before reweighting) to HERWIG 6 are given. Official 8 TeV systematic uncertainties were used for deriving the scale factor uncertainty. Multijet events were generated for proton-proton collisions at $\sqrt{s} = 8$ TeV.

	Loose	Medium	Tight
CSV discriminator	> 0.244	> 0.679	> 0.898
ZEUS+PDG+CLEO	+4.4%	+13.0%	+21.0%
Remaining differences	−4.5%	−10.9%	−17.1%
Relative differences	8.6%	20.8%	30.9%
Scale factor uncertainty	4.8%	6.8%	8.7%

Table 7.12.: Relative changes in misidentification probability for each CSV working point after reweighting c -quark jets generated by HERWIG 6 to ZEUS, PDG and CLEO values. Remaining differences (after reweighting) and relative differences (before reweighting) to HERWIG 6 are given. Official 8 TeV systematic uncertainties were used for deriving the scale factor uncertainty. Multijet events were generated for proton-proton collisions at $\sqrt{s} = 8$ TeV.

	Loose	Medium	Tight
CSV discriminator	> 0.244	> 0.679	> 0.898
ZEUS+PDG+CLEO	+0.4%	+3.7%	+8.3%
Remaining differences	+4.3%	+9.8%	+14.6%
Relative differences	8.0%	17.2%	23.6%
Scale factor uncertainty	4.8%	6.8%	8.7%

7. Impact of heavy flavour modelling on the misidentification of c-quark jets

It can be seen that the relative changes when doing the reweighting to measurements and the remaining differences after reweighting are larger than the uncertainty using official scale factor uncertainties, except for the loose working point. Relative differences before reweighting are even larger for all working points. These scale factor uncertainties are of similar order for PYTHIA 6 and HERWIG 6, since p_T binned scale factor uncertainties derived with PYTHIA 6 were used for both uncertainty calculation. Differences could come from different c-quark jet momentum spectra, but since the c-quark jet scale factor uncertainties are in the order of 10% for the dominant p_T region, these differences in the applied scale factors are too small to lead to differences in the scale-factor uncertainty between PYTHIA 6 and HERWIG 6. Differences and uncertainties in the simulation are not fully covered by the scale factor uncertainties.

7.6. Conclusions

Identification of b-quark jets is important for many measurements and searches done with the CMS experiment. The main background are c-quark jets, which have a high misidentification probability, since bottom hadrons and charmed hadrons have very similar properties. The CSV algorithm is the most commonly used b-quark identification algorithm for 8 TeV analyses.

In this chapter, it was studied if differences in heavy flavour modelling have an effect on the misidentification probability of c-quark jets. Multijet events were generated with PYTHIA 6 and HERWIG 6 for proton-proton collisions at $\sqrt{s} = 8$ TeV.

There are differences in modelling charmed hadron production fractions, charmed hadron decays and charmed hadron momentum fractions. These differences could affect the misidentification probability, as this probability depends on the charmed hadron type and the CSV algorithm uses track information and depends on the hadron momenta. It was shown that these differences affect the shape of the CSV discriminator of c-quark jets by reweighting c-quark jets generated with HERWIG 6 to PYTHIA 6 charmed hadron distributions. There are some remaining differences in the CSV discriminator shape which lead to differences in the misidentification probability. The relative difference of PYTHIA 6 to HERWIG 6 in misidentification probability of c-quark jets is 20.8% for the medium working point.

Reweighting charmed hadron distributions to measured ones also does not remove all differences in the discriminator shape. It reduces relative difference for the medium working point to 10.9%. The systematic uncertainty coming from scale factor uncertainties on the multijet sample is 6.8% for this working point. Differences in heavy flavour modelling are not covered for all three working points by official systematic uncertainties [124].

Reweighting charmed hadron production fractions, decays and momentum fractions would help to have less dependency on the used MC generator for the misidentification probability. This would also help, since scale factors are derived for events generated with PYTHIA. As not all differences are removed after reweighting, other sources have to be investigated, too.

Possible differences could arise from differences in charmed hadron mass modelling. This has an effect on the secondary vertex mass, which is used by the CSV algorithm. Differences in the charmed hadron life time modelling can affect the distance between primary and secondary vertex. Figure 7.20 shows the distribution of the secondary vertex mass and the flight distance of the secondary vertex for c-quark jets generated with PYTHIA 6 and HERWIG 6, respectively. There is a deviation between both MC generators for low secondary vertex masses, therefore charmed hadron masses should be studied, since this can affect the CSV discriminator shape. Differences for the secondary vertex flight time are smaller, but they could affect the CSV discriminator shape, therefore the charmed hadron life time modelling should be investigated.

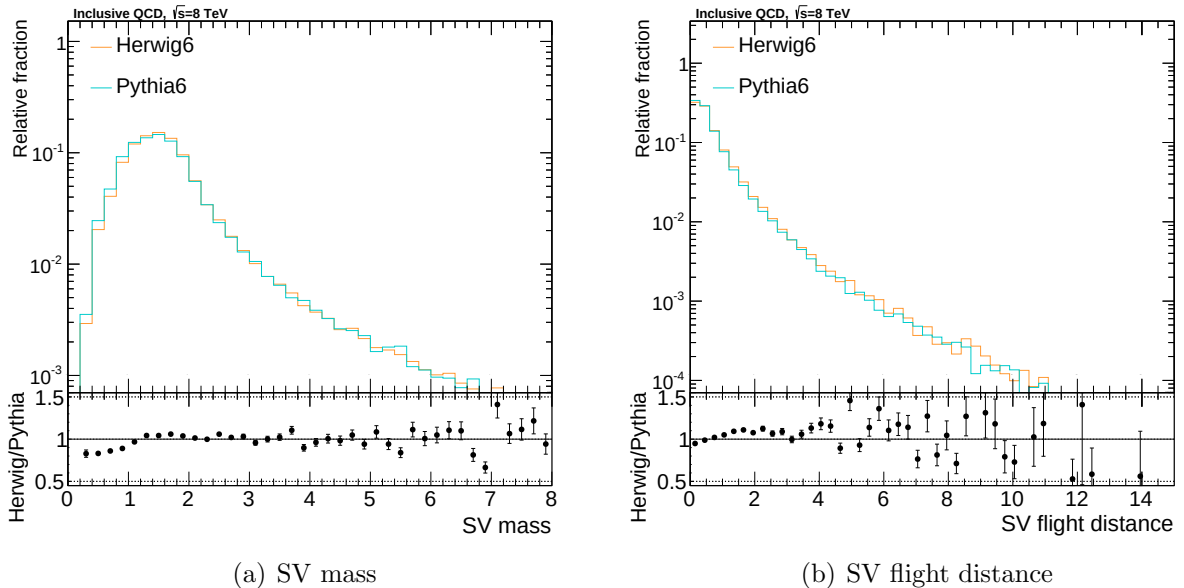


Figure 7.20.: Distributions used by the CSV algorithm for c-quark jets : (a) secondary vertex (SV) mass and (b) SV flight distance. Multijet events were generated using PYTHIA 6 (blue) and HERWIG 6 (red) for proton-proton collisions at $\sqrt{s} = 8$ TeV.

7. Impact of heavy flavour modelling on the misidentification of c-quark jets

Table 7.13.: Jet flavour fractions for jets generated with PYTHIA 6 and HERWIG 6, respectively. The ratio of the jet flavour fraction for jets generated with PYTHIA to the one generated with HERWIG is given. Jets where no parton could be assigned to are labelled as "no flavour". HERWIG 6 Multijet events were generated for proton-proton collisions at $\sqrt{s} = 8$ TeV.

Jet flavour	PYTHIA 6	HERWIG 6	PYTHIA 6/HERWIG 6
no flavour	2.7%	3.0%	0.93
g	61.7%	—	—
uds	27.4%	88.0%	0.31
c	5.7%	4.8%	1.20
b	2.5%	2.4%	1.06

Jet flavour modelling could also affect the CSV discriminator shape, when the momentum differs between both MC generators. It can have an impact on the normalisation of the c-quark jet background in comparison to light flavour and b-quark jets. Table 7.13 lists the jet flavour fractions for multijets generated with PYTHIA 6 and HERWIG 6. There are categories for gluon jets (g), light flavour jets (uds), c-quark jets (c), b-quark jets (b), and a category for jets where no parton could be assigned to (no flavour). There are no gluon jets in samples generated with HERWIG 6 as gluons are forced to split into quark-antiquark pairs at the end of the parton shower, see Chapter 5. Multijet events generated with PYTHIA 6 have a higher fraction of c-quark jets than events generated with HERWIG 6. This leads to a different amount of c-quark jet background depending on which MC generator was used, especially because the b-quark jet fraction is roughly the same.

There are differences in heavy flavour modelling between PYTHIA 6 and HERWIG 6 which affect the misidentification probability for b-quark jet identification. These differences can also affect analyses with c-quark jets, for example searches for charged Higgs bosons which decay into a s-quark and a c-quark ($H^\pm \rightarrow cs$), or into a c-quark and a b-quark ($H^\pm \rightarrow cb$).

Identification of c-quark jets would make it easier to measure scale factors for b-quark jet identification. It could help to reduce background of charged Higgs boson searches and other analyses with c-quark jets in the final state.

Charmed hadron modelling has to be improved for such analyses and c-quark jet identification algorithms, as it was shown that it already affects existing algorithms. Since there were no charmed hadron measurements done at the LHC relevant for b-quark jet identification at time of this study, these observables should be measured at LHC energies.

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

Lepton-flavour violating (LFV) decays of the Higgs boson can be introduced by flavour-violating Yukawa couplings. In this chapter, a search for lepton-flavour violating (LFV) decays of the Higgs boson to an electron and a tau, denoted $e\tau$, which further decays to a muon and neutrinos is presented. This decay channel is denoted $e\tau_\mu$. The search was performed for proton-proton collisions at $\sqrt{s} = 8$ TeV and $\sqrt{s} = 13$ TeV. Results of the search at $\sqrt{s} = 8$ TeV were already published together with the results for the hadronically decaying tau, denoted $e\tau_h$ channel, and the search for LFV decays of the Higgs boson to an electron and a muon, denoted $e\mu$ channel. All three searches were done at $\sqrt{s} = 8$ TeV and the results were published in [5]. As the analyses, done at $\sqrt{s} = 8$ TeV and $\sqrt{s} = 13$ TeV, are based on very similar techniques, this chapter focus on the $\sqrt{s} = 13$ TeV search. Differences between those analyses are explained when appropriate later.

First, the signal process and the mass reconstruction are described. Then, background processes and the preselection of events are discussed, followed by a description of the background modelling. A difference between both analyses is the background modelling of the two dominant backgrounds. All background methods are discussed in this thesis and it is explained why different background methods were used. Control regions for important background processes are also discussed. Afterwards, the profile likelihood method is described, which is used to fit the background and signal model to the collision data. It is also used to derive limits on the branching fraction of LFV decays of the Higgs boson. Systematic uncertainties, which are nuisance parameters of the fit, are introduced for the analysis at $\sqrt{s} = 13$ TeV. Variables which have discriminating power to separate signal from background processes are presented. Then, the search at $\sqrt{s} = 8$ TeV is summarised. Afterwards, the optimisation of the event selection for $\sqrt{s} = 13$ TeV using a self-developed algorithm is discussed. Then, the expected upper limits on the branching fraction of LFV decays of the Higgs boson are presented and the event selection for $\sqrt{s} = 13$ TeV is derived. Finally, the results are presented for the CMS collision data collected in 2015 at $\sqrt{s} = 13$ TeV corresponding to an integrated luminosity of 2.3 fb^{-1} and they are compared to the results of the search at $\sqrt{s} = 8$ TeV [5], which was performed using CMS collision data collected in 2012 corresponding to an integrated luminosity of 19.7 fb^{-1} .

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

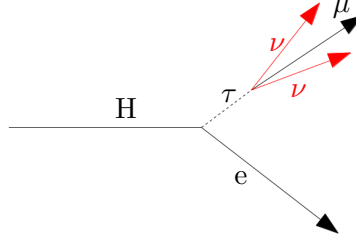


Figure 8.1.: Feynman diagram of the leptonic lepton-flavour violating decay of the Higgs boson to $e\tau$ with an electron-muon plus $\cancel{E}_T$ final state.

8.1. Signal process

If the observed Higgs-boson couples to the leptons according to their mass, see Section 3.3, the branching fraction of LFV $H \rightarrow e\tau$ decays is the second largest one of the possible three LFV Higgs-boson decays. Indirect constraints on the $H \rightarrow e\tau$ branching fraction are in the same order as for $H \rightarrow \mu\tau$. Since both signatures are similar, an improvement of these upper limits by a direct search using the CMS data is expected.

The leptonic decay of the tau to a muon has a smaller branching fraction than the hadronic tau decay, but it leads to a cleaner detector signature. Muons can be better identified within the CMS experiment than electrons or hadronically decaying taus leading to a smaller fraction of misidentified leptons.

Figure 8.1 shows the corresponding Feynman diagram for the LFV $H \rightarrow e\tau_\mu$ decay. The final state includes an isolated electron and an isolated muon of opposite charge. Neutrinos of the tau lead to missing transverse energy. This final state is very similar to SM decays of the Higgs boson to taus which decay further to an electron, a muon, and neutrinos. Main differences between both signatures are the prompt electron for LFV decays and the electron from tau decays for SM decays.

Events with SM and LFV Higgs-boson decays were generated for proton-proton collisions at $\sqrt{s} = 13$ TeV and a bunch spacing of 25 ns using POWHEG for the matrix elements and PYTHIA 8 for decay and hadronisation. The CMS Underlying Event Tune for PYTHIA 8 (CUETP8M1) [118] and a pileup profile derived from 2015 observed data were used. Detector geometry and detector response were simulated with GEANT4. Table B.1 from Appendix B gives an overview of the MC samples used in this analysis. Cross sections for the processed can be found in Table B.2 from Appendix B.

The SM decay of the Higgs boson to taus is a background of a search for LFV decays of a Higgs boson to $e\tau_\mu$. Figure 8.2 shows the difference in the shape of the kinematic distributions between LFV and SM Higgs-boson decays. Variable distributions are shown for (a) electron momentum, (b) muon momentum, (c) $\Delta\phi(\mu, \vec{E}_T)$ and (d) $\Delta\phi(e, \vec{E}_T)$, the difference in ϕ between the muon or electron and the missing-energy vector. The LFV Higgs-boson decay has a harder electron momentum spectrum, as the electron is a prompt lepton and not coming from a tau decay as in the SM case. Since in both cases the muon is coming from a tau decay, its momentum spectrum is soft and very similar for LFV and SM Higgs-boson decays. But there are differences between LFV and SM decays in $\Delta\phi(\mu, \vec{E}_T)$. In LFV $H \rightarrow e\tau_\mu$ decays, neutrinos are only originating from one tau decay. In addition, the mass difference between the Higgs boson and the tau is very large, therefore $\vec{E}_T$ tends to be closer to the muon in the LFV case. The $H \rightarrow \tau\tau$ background process can be suppressed by requirements on the electron momentum and the angle between the direction of the muon and of $\vec{E}_T$.

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

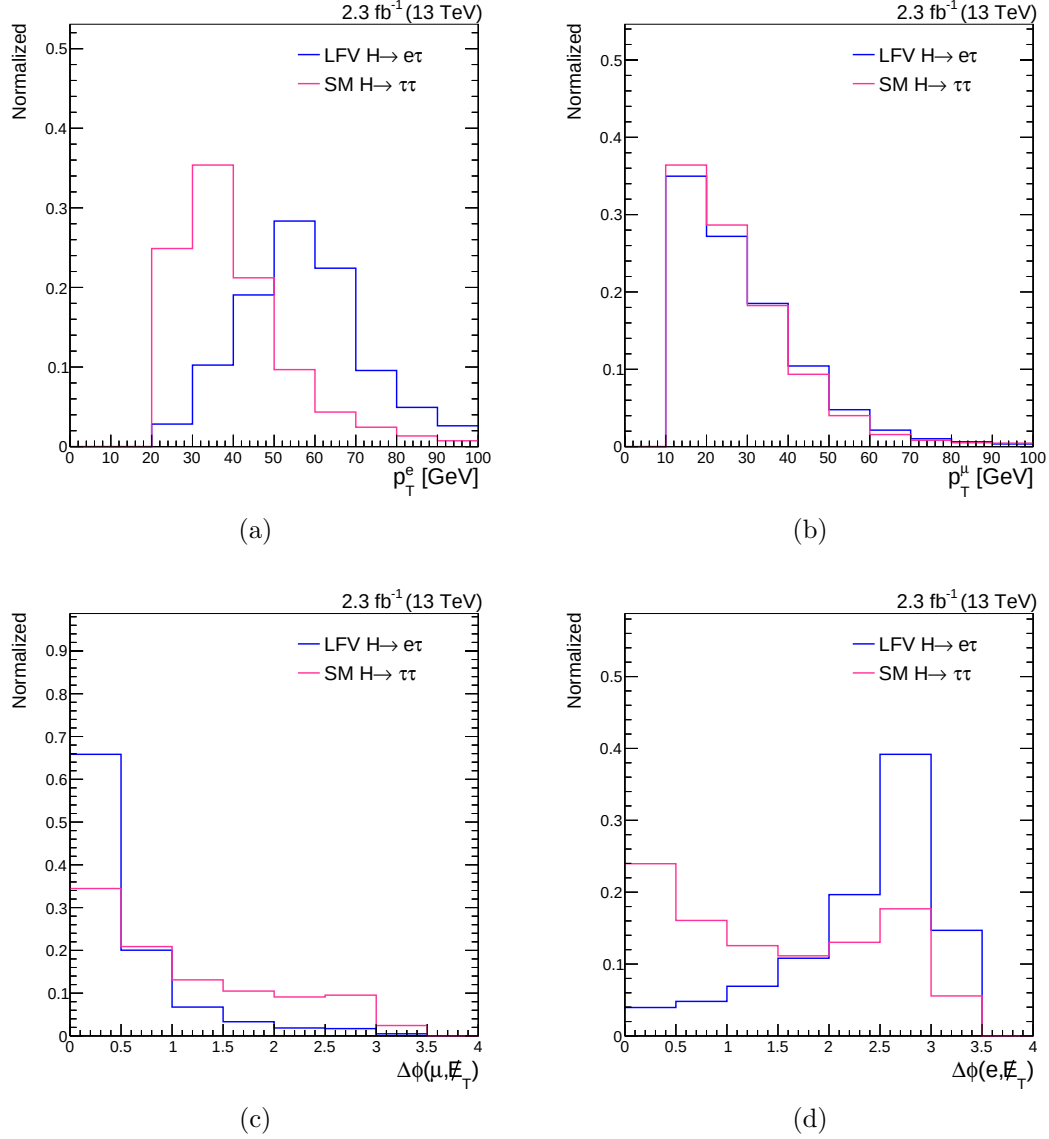


Figure 8.2.: Shape of variable distributions for decays of the Higgs boson to $e\tau_\mu$ (blue) or $\tau_e\tau_\mu$ (pink), respectively. The distributions have been normalised by their total number of events. They are shown for (a) electron transverse momentum, (b) muon transverse momentum, (c) $\Delta\phi(\mu, \vec{E}_T)$, and (d) $\Delta\phi(e, \vec{E}_T)$.

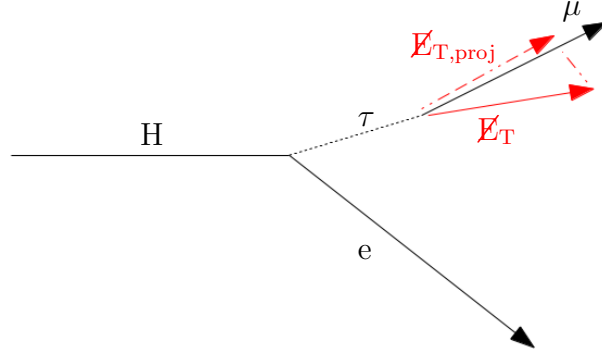


Figure 8.3.: Estimation of the neutrino momentum $\vec{E}_{T,\text{proj}}$ by using the component of the missing transverse energy $\vec{E}_T$ which is collinear to the muon flying-direction in the transverse plane. Shown are the Feynman diagram of the LFV Higgs-boson decay (black), the missing transverse energy (red vector) and the projected missing transverse energy $\vec{E}_{T,\text{proj}}$ (red, dashed vector).

8.2. Mass reconstruction

The mass reconstruction is very important for the search for LFV decays of the Higgs boson to $e\tau_\mu$. Since only the transverse energy of the proton-proton collisions is known, $\vec{E}_T$ and related quantities are calculated in the transverse plane. As the tau is decaying to a muon and neutrinos, the visible mass M_{vis} , which is defined as the invariant mass of the electron and the muon, is smaller than the mass of the Higgs boson. Neutrinos lead to missing transverse energy, $\vec{E}_T$, in the detector. As described in Section 6.7, not only genuine $\vec{E}_T$ contributes to the reconstructed missing transverse energy, but there are also contributions from jet mismeasurements.

Since jet mismeasurements can lead to large artificial $\vec{E}_T$ contributions, missing transverse energy alone is not well suited for reconstruction of the sum of neutrino momenta, but the event topology of the signal can be used to improve the mass reconstruction. Taus have a much smaller mass than the Higgs boson, therefore their decay products are boosted, consequently they are collinear to the flight-direction of the tau. Thus, the neutrinos are collinear to the flying direction of the visible tau decay products, which is the muon for the $e\tau_\mu$ channel. The neutrino momentum can be estimated by the missing transverse energy momentum which is collinear to the muon, $\vec{E}_{T,\text{proj}}$:

$$\vec{E}_{T,\text{proj}} = \frac{\vec{E}_T \cdot \vec{p}_T^{\text{vis}}}{p_T^{\text{vis}}} \approx p_T^\nu. \quad (8.1)$$

Figure 8.3 visualises the projected missing transverse energy $\vec{E}_{T,\text{proj}}$ as red, dashed vector.

The collinear mass, M_{col} , is defined as invariant mass of the electron, visible tau decay products (the muon) and the neutrinos ($\vec{E}_{T,\text{proj}}$). Using the collinear approximation, the mass of the Higgs boson can be expressed in terms of the visible mass and the visible

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

fraction of the tau decay, where the visible fraction of the tau decay products, x_τ^{vis} , is defined as:

$$x_\tau^{\text{vis}} = p_T^{\tau^{\text{vis}}} / (p_T^{\tau^{\text{vis}}} + \cancel{E}_{T,\text{proj}}). \quad (8.2)$$

In Appendix A.1, the following expression for the collinear mass is derived from the invariant mass of the Higgs boson:

$$M_{\text{col}} = M_{\text{vis}} / \sqrt{x_\tau^{\text{vis}}}. \quad (8.3)$$

Figure 8.4 compares the visible mass and the collinear mass for LFV decays of the Higgs boson to $e\tau_\mu$ in events with no reconstructed jet with a momentum of $p_T > 30$ GeV, which is the transverse momentum requirement used in all searches for LFV decays of the Higgs boson at the CMS experiment [5, 17, 125]. The mass resolution of events with one or two of such jets is shown in Figure A.1 from Appendix A. The mean of the visible mass is below the Higgs-boson mass, while the mean of the collinear mass is very close to the mass of the Higgs boson. In addition, the resolution of the collinear mass is better than the resolution of the visible mass by a factor of 1.6. The behaviour between both mass quantities is similar for events with one or two jets.

For these reasons the collinear approximation, used for the reconstruction of the sum of neutrino momenta, makes it possible to calculate the invariant mass of the LFV Higgs-boson decay products. As the collinear mass distribution, M_{col} , has a mean value close to the Higgs-boson mass of 125 GeV and a better resolution than the visible mass, it is used to define the signal region as $100 \text{ GeV} < M_{\text{col}} < 150 \text{ GeV}$.

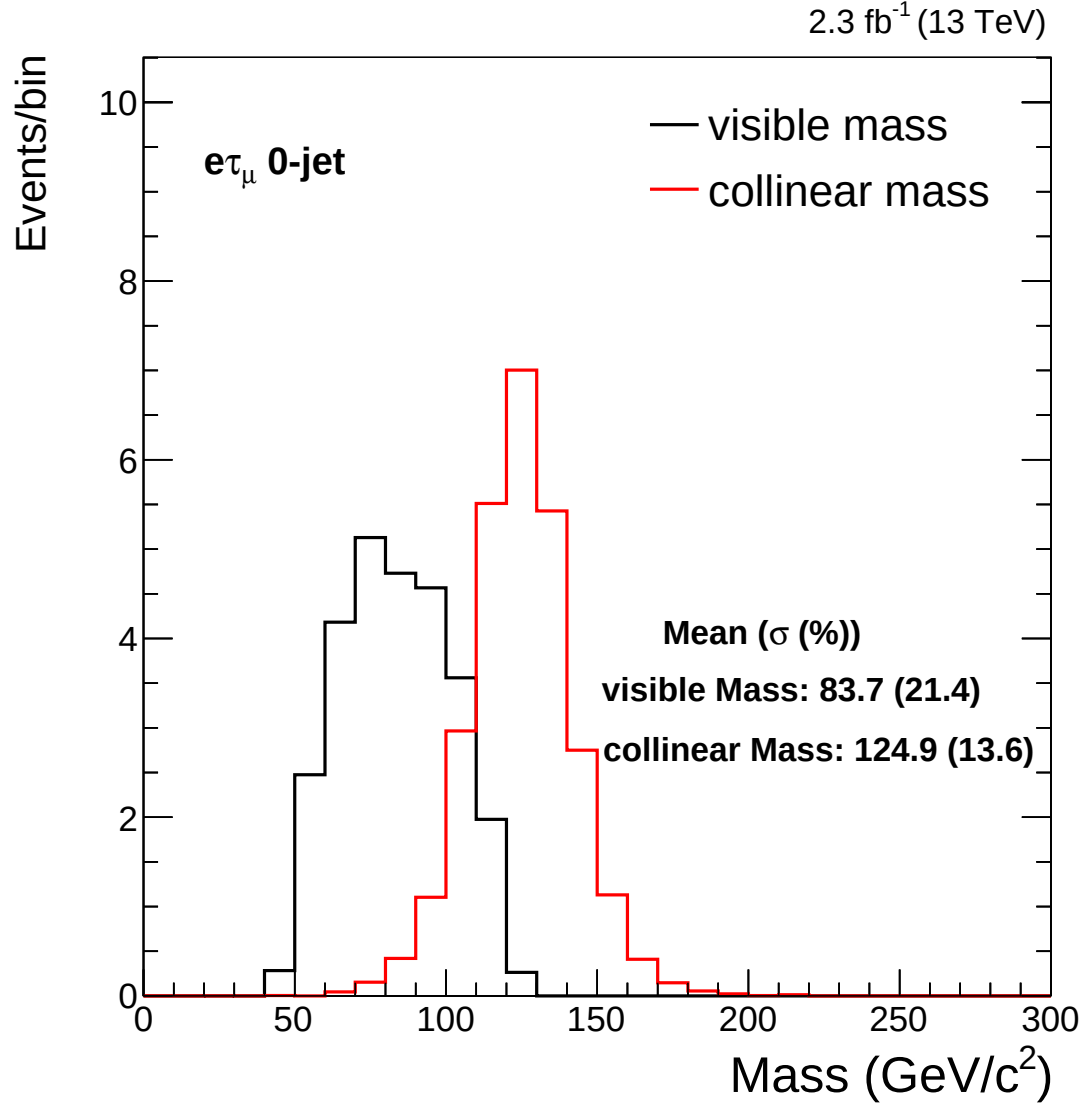


Figure 8.4.: Visible mass (black line), and collinear mass (red line) are shown for LFV decays of the Higgs boson to $e\tau_\mu$. Events shown in the plot have no reconstructed jet with $p_T > 30$ GeV.

8.3. Background processes

Different background processes leave a signature in the detector similar to LFV decays of the Higgs boson to $e\tau_\mu$. Figure 8.5 shows the Feynman diagrams of background processes which have a $e\mu + \cancel{E}_T$ signature. The SM $H \rightarrow \tau_e\tau_\mu$ background has a signature, which is very close to the LFV case, except that the electron originates from a tau decay, therefore it is softer. Since neutrinos from both tau decays contribute to $\cancel{E}_T$ and the collinear mass approximation uses only that component of $\cancel{E}_T$ which is parallel to the muon direction, the resolution of the collinear mass is by a factor 1.5 – 1.7 worse for the SM $H \rightarrow \tau_e\tau_\mu$ case than for LFV Higgs decay to $e\tau_\mu$, Figure A.2 from Appendix A. The process $Z \rightarrow \tau_e\tau_\mu$ has a similar behaviour. Additionally, the peak of the collinear mass is shifted to lower masses with respect to the Higgs mass. Due to the large production cross section of Z bosons compared to the Higgs-boson production cross-section, $Z \rightarrow \tau_e\tau_\mu$ is the dominant background process for events which have a low jet multiplicity. A higher requirement on the electron transverse momentum is expected to reduce the SM Higgs boson and the $Z \rightarrow \tau_e\tau_\mu$ background contributions.

Another important background for low jet multiplicities is the diboson background which is dominated by the WW process. Both W bosons can decay to a lepton and a neutrino. The resulting missing transverse energy is not collinear to the muon direction, therefore the collinear-mass resolution of this process is poor. The WZ and the ZZ process have lower cross sections and at least one Z boson is required to decay to taus with at least one tau decaying leptonically. These processes can have further leptons, therefore a veto on further leptons will be applied to reduce these background processes. As the diboson background consists of processes which lead to a higher invariant mass than the Higgs mass, the collinear mass is shifted to higher values than for the LFV Higgs boson case.

Another important background, especially for higher jet multiplicity events, is the top background. Both, the di-leptonic $t\bar{t}$ process and the tW process contribute to the top background. This background leads to b-quark jets and it can be reduced by requiring that none of the jets is identified as a b-quark jet. Neutrinos of both W-boson decays and jet mismeasurements contribute to $\cancel{E}_T$, leading to a poor collinear mass resolution. This background also peaks at higher collinear mass values due to the large top mass.

A minor background is $W\gamma/\gamma^*$. This background has a low cross-section times branching-fraction value.

Misidentified leptons are a very important background. This background is dominated by the W+jets process. The W boson can decay leptonically and another Particle-flow (PF) candidate, Section 6.5, such as a lepton within a jet or a charged hadron, can be misidentified as a lepton. Leptons within jets have a significant energy flow around their trajectory. Tighter requirements on the lepton isolation can reduce the misidentification

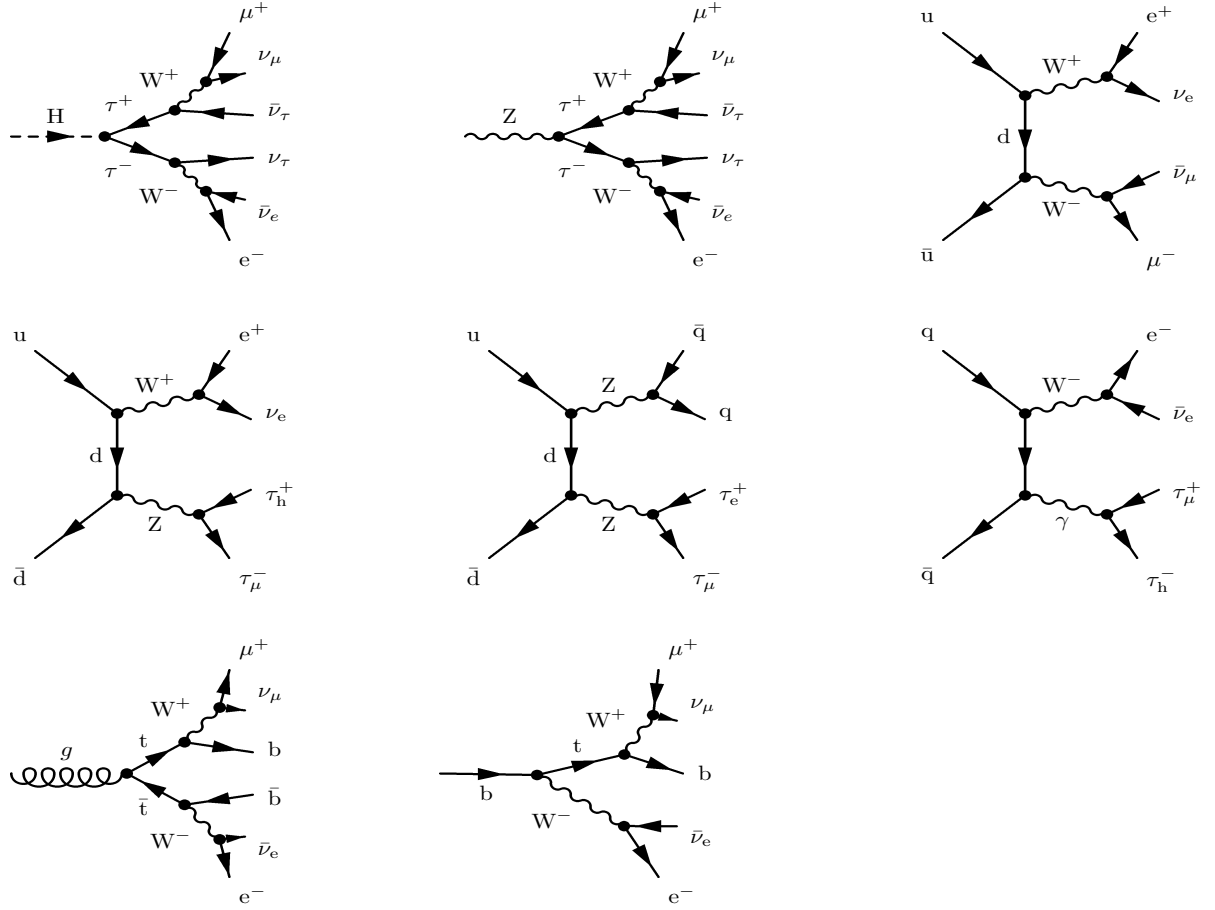


Figure 8.5.: Feynman diagrams of background processes to LFV Higgs boson decays to $e\tau_\mu$ which have a $e\mu + \cancel{E}_T$ signature: (a) $H \rightarrow \tau_e \tau_\mu$, (b) $Z \rightarrow \tau_e \tau_\mu$, (c) WW , (d) WZ , (e) ZZ , (f) $W\gamma/\gamma^*$, (g) $t\bar{t}$, and (h) $tW/\bar{t}W$.

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

Table 8.1.: Overview of properties of the background processes - The properties of the background processes are compared to the ones of the signal process (highlighted in grey), which are the lepton multiplicity N_l (without hadronic taus), lepton charges, jet multiplicity N_{Jet} , and the b-quark jet multiplicity. A further property is if the processes can have an electron-muon signature requiring both leptons to be prompt or from a tau decay. If this is not the case, such a signature can only arise if either one of both leptons was misidentified or if both leptons were misidentified.

Process	N_l	2 prompt leptons/ τ_l	$e^\pm\mu^\mp$	$e^\pm\mu^\pm$	N_{Jet}	b-quarks
$H \rightarrow e\tau_\mu$	2	✓	✓	-	0/2	-
$H \rightarrow \tau_e\tau_\mu$	2	✓	✓	-	0/2	-
$Z \rightarrow \tau_e\tau_\mu$	2	✓	✓	-	0	-
W+jets	1	-	✓	✓	0	-
WW	2	✓	✓	-	0	-
WZ	2-3	✓	✓	✓	0-1	-
ZZ	2-4	✓	✓	✓	0-2	-
$t\bar{t}$ (di-leptonic)	2	✓	✓	-	2	2
tW (di-leptonic)	2	✓	✓	-	1	1

probability. Jets consist of hadrons, which can be only described with phenomenological models, therefore processes with misidentified leptons have the highest systematic uncertainty. Other processes, which contribute to the misidentified-lepton background, are diboson processes and processes including top quarks. Bottom hadrons have large branching fractions to leptons, which results in a higher misidentification probability. The $(Z \rightarrow ll) + \text{jets}$ processes, with $l = e, \mu, \tau$ also contribute to the misidentified-lepton background. Another type of events that can lead to two misidentified leptons are QCD multijet events, which are negligible due to the small probability for two leptons to be misidentified.

Table 8.1 gives a summary of the background properties discussed. Due to the large cross section, $Z \rightarrow \tau_e\tau_\mu$ is the dominant background process for low jet multiplicities. At higher jet multiplicities, the top background becomes more important. The misidentified-lepton background, which is dominated by the W+jets process, is important, especially for low jet multiplicities, since it has a high systematic uncertainty. Diboson processes, such as WW, contribute, too. A good preselection is necessary to reduce the SM background contributions.

Detailed MC simulations of all background sources are employed for the background modelling. Events were generated with PYTHIA 8 for proton-proton collisions at $\sqrt{s} = 13$ TeV and a bunch spacing of 25 ns. The CMS Underlying Event Tune for PYTHIA 8 (CUETP8M1) [118] and a pileup profile derived from 2015 collision data were used. Detector geometry and detector response were simulated with GEANT4. Table B.1 from Appendix B gives an overview of the MC samples used in this analysis. Cross sections for the processed can be found in Table B.2 from Appendix B.

8.4. Preselection of events

In this section, the preselection of events is discussed for the search at $\sqrt{s} = 13$ TeV. The aim of the preselection is to select clean events with a high efficiency for LFV Higgs boson decays to $e\tau$, while reducing backgrounds on the other hand. A data sample is used which was collected with an unscaled trigger. Physics object definitions are used, described in Chapter 6, which reduce the amount of leptons within jets. Since observed and simulated properties do not have to agree, corrections are applied, which were derived from other processes in collision data. Finally, further requirements are discussed and events are classified according to their jet multiplicity.

8.4.1. Recorded data sample

The data sample used in this search was collected in proton-proton collisions at $\sqrt{s} = 13$ TeV with a bunch spacing of 25 ns. It was recorded with the CMS detector in 2015 and corresponds to an integrated luminosity of 2.32 fb^{-1} .

As LFV decays of the Higgs boson to the $e\tau$ has a p_T -spectrum which peaks at $p_T \sim 60$ GeV, Figure 8.2(a), events are required to have been recorded with a single electron trigger, `HLT_Ele23_WPLoose_Gsf`, with a requirement on the online supercluster energy of $E_T > 23$ GeV. There is no requirement on the pseudorapidity, therefore reconstructed electrons have a pseudorapidity of $|\eta^e| < 2.5$, as described in Section 6.2. Loose identification and isolation criteria are also applied by the trigger. The identification criteria are based on the transverse profile of the clustered energy in the ECAL, the ratio of the HCAL energy to the total supercluster energy, and the degree of association between the track and the ECAL cluster [55]. Isolation is based on the amount of energy surrounding the electron in the tracker, ECAL, and HCAL. This trigger was not prescaled in 2015, but it was used a short time after collecting proton-proton collisions with a 25 GeV bunch spacing. This affects only a very small luminosity fraction and runs before the start of this trigger usage are not considered. Other unscaled single electron triggers, with different p_T requirements and with a requirement on the pseudorapidity, are available, but they have lower efficiencies with respect to simulated LFV $H \rightarrow e\tau_\mu$ events.

A muon-electron trigger with a higher requirement on the electron p_T than the muon p_T has not been considered because such a trigger has also a requirement on the muon isolation. Since the p_T -spectrum of jets is steeply falling as function of p_T , the misidentification rate of PF candidates as leptons depends on their transverse momentum. These PF candidates carry just a fraction of the quark energy. Thus, the misidentification rate is higher at smaller p_T values. Muons originating from tau decays have a softer p_T -spectrum than the electrons, Figure 8.2(b). Thus, the misidentification rate of a PF candidate as a muon is higher than for the electrons. The fraction of misidentified leptons is lower for tighter

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

isolation requirements. A misidentified-lepton side-band can be defined by loosening the isolation criteria. This isolation side-band can be employed for the background modelling of the misidentified lepton background, described in Section 8.5.2. This is less complicated if there are no requirements applied on the muon by the trigger, as the possible isolation side-band is not limited by the trigger and there is no preselection on the muon due to the trigger.

8.4.2. Physics objects

In this section, physics objects, used in the search for LFV decays of the Higgs boson to $e\tau_\mu$, are defined. Electrons, muons, and missing transverse energy, due to the neutrinos, are part of the LFV Higgs-decay signature. Jets are important for employing the jet multiplicity of different Higgs-boson production-mechanisms. Identification of b-quark jets can be used to reduce the background contribution coming from events with top-quarks.

Electrons Electrons are reconstructed using the techniques described in Section 6.2. As prompt electrons have to be separated from background sources, requirements on ID observables are applied. Either a cut-based approach or an MVA technique can be applied. This analysis uses the MVA technique for electron identification. A BDT can be trained with all electrons, for the non-triggering MVA ID, or it can be trained with electrons which pass typical electron-trigger requirements, for the triggering MVA ID. The triggering MVA ID has same advantages for electrons recorded with electron triggers. Different identification and isolation observables and criteria are used by electron identification and electron triggers. Furthermore, the HLT menu in MC simulation is not necessarily the same as for collision data, therefore electrons have to pass typical HLT criteria, loose identification and isolation criteria, to create a common baseline for simulated and with an electron trigger recorded events. Applying the ID on electrons which do not pass the triggering selection will result in undefined behaviour, therefore these criteria have to be applied for electrons which pass the triggering MVA ID. Since events used in this analysis have been recorded with a single electron trigger, ID criteria which have been optimised for such electrons have the best performance. Thus, a triggering MVA ID is applied on electrons. Additionally, this makes the data analysis less dependent on trigger MC simulation. The working point *wp 80* is used which has an identification efficiency of 80% for prompt electrons, which corresponds to the efficiency of the medium cut-based ID that is designed for measurements involving W or Z bosons. Table 8.2 lists the triggering MVA selection. The criteria differ between electrons reconstructed in the barrel (EB) and electrons which have been reconstructed in the endcap (EE). Important identification observables, such as the width of the supercluster in η , are used. Additionally, there are some requirements on the electron momentum and the PF-cluster isolation in the ECAL,

Table 8.2.: Triggering MVA selection for 2015 data with a 25 ns bunch spacing. Electron p_T , width of the SC in η ($5 \times 5 \sigma_{i\eta i\eta}$), spatial matching between track and SC ($|\Delta\eta_{in}|$ and $|\Delta\phi_{in}|$), the ratio of HCAL energy and SC energy (H/E), and the relative isolation in the ECAL, HCAL, and Tracker.

Variable names	Barrel (EB)	Endcap (EE)
p_T	$> 15 \text{ GeV}$	$> 15 \text{ GeV}$
$5 \times 5 \sigma_{i\eta i\eta}$	< 0.012	< 0.033
$ \Delta\eta_{in} $	< 0.0095	-
$ \Delta\phi_{in} $	< 0.065	-
H/E	< 0.09	< 0.09
ECAL PF-cluster isolation/ p_T	< 0.37	< 0.45
HCAL PF-cluster isolation/ p_T	< 0.25	0.28
Tracker isolation/ p_T	< 0.18	0.18

HCAL, and the Tracker. Table 8.3 gives an overview of the requirements on the MVA discriminator, which depend on the supercluster η . The requirements differ for electrons which have been reconstructed in the barrel or in the endcap. Furthermore, the barrel region is split into two η regions. Both, the triggering MVA selection and the requirements on the MVA discriminator, have to be applied for the triggering MVA ID.

Table 8.3.: Requirements on the MVA discriminator of the triggering MVA ID wp 80. The identification criteria depend on the supercluster η .

Category	MVA discriminator
Barrel ($ \eta < 0.8$)	> 0.988153
Barrel ($ \eta > 0.8$)	> 0.967910
Endcap	> 0.841729

Electron isolation can be used to further separate prompt electrons from electrons within jets or jets misidentified as electrons. The relative PF isolation I_{PF}^{rel} , described in Section 6.2.5, is used with a cone radius of $R = 0.3$. Additional energy due to pileup is subtracted using rho-effective area corrections. The effective area depends on the absolute value of the supercluster η and is between 0.14 and 0.27. A tight requirement of $I_{PF}^{rel} < 0.1$ is used. Finally, requirements on the electron kinematics are applied. Electrons are required to have a transverse momentum $p_T^e > 20 \text{ GeV}$ and a pseudorapidity of $|\eta| < 2.5$, excluding the transition region between EB and EE $1.4442 < |\eta| < 1.566$.

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

Muons Muons are reconstructed using the techniques described in Section 6.3. The *Medium Muon ID* newly introduced in LHC Run II is designed to be efficient for prompt muons and for muons from heavy-quark decays. It was designed to have less misidentified muons than the loose ID and to be more efficient than the tight ID.

Table 8.4 shows the identification criteria for the Medium Muon ID. PF Muons are selected if they are reconstructed as a global muon or as a tracker muon. They have to satisfy a requirement on the fraction of valid tracker hits. Additionally, they have to fulfil quality criteria to be classified as a good global muon or they have to pass the tight segment compatibility criterion.

Table 8.4.: Criteria of the Medium Muon ID- Muons which pass the medium Muon ID also pass the loose Muon ID. They also have a minimum fraction of valid tracker hits. Additionally, they have to either pass the good global muon requirements, or they have to satisfy the tight segment compatibility.

Loose Muon ID	PF muon
	global or tracker muon
Fraction of valid tracker hits	> 0.8
1) Good global muon	global muon
	normalised χ^2 of global track $\chi^2 < 3$
	tracker-standalone position match $\chi^2 < 12$
	Kink finder track kink < 20
	segment compatibility > 0.303
2) Tight segment compatibility	segment compatibility > 0.451

The background contributions from heavy-flavour quarks and hadrons decaying into muons is reduced by a requirement on the relative PF isolation. A cone of radius $R = 0.4$ is used to estimate the energy flow around the muon. Contributions in the energy sum due to pileup are subtracted using $\Delta\beta$ corrections [110]. The correction factor β is estimated by multiplying the energy deposit of charged particles which are not associated with the primary vertex by a factor of 0.5. It has been empirically found that the total energy for neutral hadrons and photons originating from pileup combined is, on average, about one half of the total energy from charged hadrons, which do not come from the chosen primary vertex. A tight requirement is applied on the relative PF isolation $I_{\text{PF}}^{\text{rel}} < 0.15$, which has an efficiency of 95.7%.

Finally, kinematic requirements are applied. Muons with a transverse momentum $p_{\text{T}}^\mu > 10$ GeV and a pseudorapidity $|\eta^\mu| < 2.4$, which corresponds to the coverage of the muon system, are selected.

Table 8.5.: Loose Jet ID- Requirements are listed depending on the jet η . EM stands for electromagnetic and refers to either photons (neutral EM), or charged leptons (charged EM).

For $ \eta \leq 3$	
Neutral hadron fraction	< 0.99
Neutral EM fraction	< 0.99
Number of constituents	> 1
For $ \eta \leq 2.4$ in addition apply	
Charged hadron fraction	> 0
Charged multiplicity	> 0
Charged EM fraction	< 0.99
For $ \eta > 3$	
Neutral EM fraction	< 0.90
Number of neutral particles	> 10

Jets The PF jets are used which have been reconstructed, as described in Section 6.6, using the anti- k_t algorithm and a jet distance parameter $R = 0.4$. A *Loose Jet ID* is applied which was constructed in order to reject fake jets, badly reconstructed and noise jets, while retaining an efficiency of 98 – 99% for genuine jets. Correlations and the discriminating power of all PF jet variables as well as variables for hadronic and electromagnetic calorimeter noise rejection are exploited.

Table 8.5 lists the loose Jet ID criteria, which depend on the jet η . Track information is only available for $|\eta| \leq 2.5$ and the muon system covers the region $|\eta| \leq 2.4$, therefore observables which make use of tracking information are restricted to this jet- η region. ECAL and HCAL cover the region $|\eta| < 3$, while the HF covers the region $3 < |\eta| < 5$. A kinematic selection is applied requiring $p_T > 30$ GeV for the jet transverse momentum and $|\eta| < 4.7$ for the jet pseudorapidity. The criterion on η is smaller than the covered region to avoid boundary effects of the jet reconstruction.

Missing transverse energy Reconstruction of missing transverse energy $\cancel{E}_T$ is described in Section 6.7. Muons deposit a negligible amount of their energy in the calorimeters. As muons and electrons are a part of the signature of LFV decays of the Higgs boson to $e\tau_\mu$, tracking information and information from the muon system has to be included in the $\cancel{E}_T$ reconstruction. PF $\cancel{E}_T$ is used, since it combines the information of all subdetectors. Filters are applied to remove events with large artificial $\cancel{E}_T$, for example resulting from detector malfunctions.

Identification of b-quark jets Jets originating from b-quarks are identified using the combined secondary vertex CSVv2 algorithm as presented in Section 6.8. The medium working point is used, which has a misidentification efficiency of 1% for light-flavour jets. Jets with a CSVv2-discriminator value larger than 0.8 are identified as a b-quark jet.

8.4.3. Corrections

Observed and simulated events can differ with respect to the number of pileup events, certain efficiencies for reconstruction and identification, and the trigger efficiency. A pileup reweighting is applied on simulated events to correct the number of interactions to measurements. Reconstruction, identification, isolation criteria, and trigger efficiencies are corrected for simulated events with scale factors. A tag-and-probe method is used for measuring the efficiencies. This method is described in Section 6.4. The jet energy has to be corrected for observed and simulated jets by applying jet energy corrections. These JEC have to be propagated to the missing transverse energy. Scale factors for b-quark jet identification are discussed, too. However, they are not applied as only the normalisation is affected. Furthermore, the normalisation of the background models are fitted to collision data.

Pileup reweighting The number of additional pileup interactions can be estimated either by counting the number of reconstructed primary vertices taking the vertex reconstruction efficiency into account or by measuring the instantaneous luminosity for each bunch crossing and multiply it by the inelastic cross section, as shown in Section 6.6.3. The luminosity-based pileup estimate is used for pileup calculations and the vertex comparison is used as a cross check.

A single JSON format file stores the pileup information and allows to calculate the observed-data pileup-distribution for the luminosity sections used in a certain analysis. A minimum bias cross section of $\sigma_{\text{MB}} = 69 \text{ mb}$ with an uncertainty of 5% is used, which corresponds to the best fit value and differs from the former recommended value of $\sigma_{\text{MB}} = 80 \text{ mb}$.

The number of true interactions is also stored for events from MC simulation. Distributions of the number of true interactions can be compared between the one expected in collision data and the one for simulation. As the simulated distribution usually differs from the observed one, a pileup reweighting is applied. Simulated events are reweighted in a way that their distribution matches the one expected for collision data after reweighting. For the simulated distribution of interaction multiplicity, a $DY \rightarrow \ell\ell$ MC sample generated with MadGraph and PYTHIA 8, Table B.1 from Appendix B, is used without applying a trigger or an event selection.

Figure 8.6 shows the number of interactions expected in collision data and the one from simulated events, (a) before pileup reweighting and (b) after pileup reweighting. Both distributions match after applying pileup reweighting.

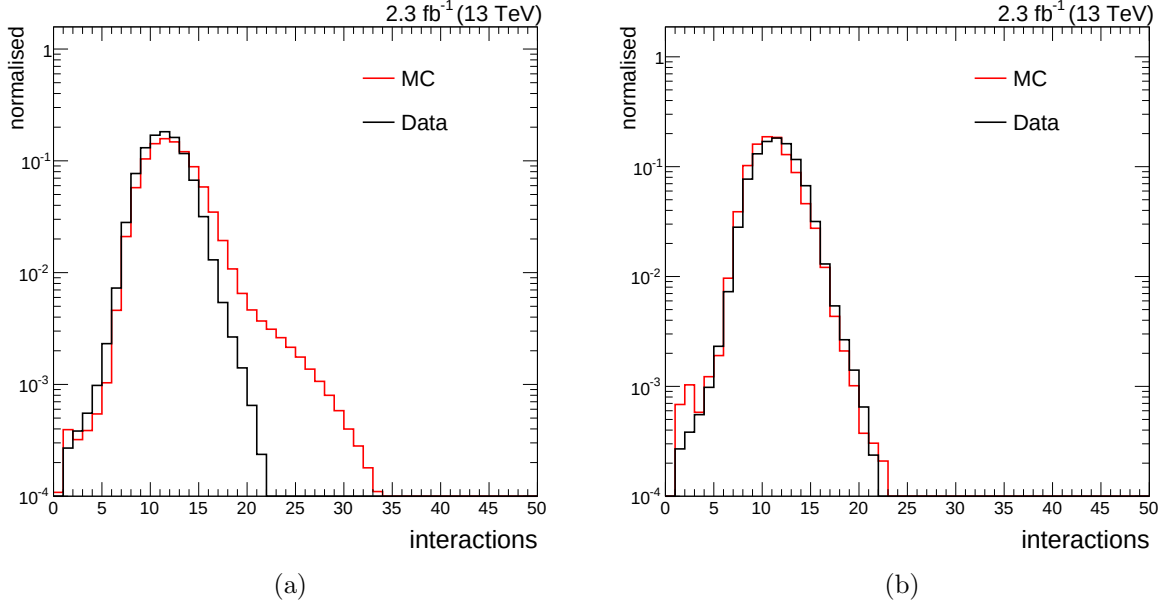


Figure 8.6.: Number of interactions for collision data and MC simulation, (a) before and (b) after reweighting simulated events to match the distribution expected in collision data.

Electron scale factors Each step of the electron reconstruction and selection has its own efficiency. These efficiencies can be different for electrons in collision data and in simulation, therefore scale factors are estimated to correct for these differences. The total efficiency for electron reconstruction and identification in this analysis is factorised in the following steps:

$$\varepsilon^e = \varepsilon_{\text{Reco}}^e \cdot \varepsilon_{\text{ID}|\text{Reco}}^e \cdot \varepsilon_{\text{Trigger}|\text{ID}}^e \cdot \varepsilon_{\text{Iso}|\text{Trigger}}^e, \quad (8.4)$$

where ε^e is the total efficiency, $\varepsilon_{\text{Reco}}^e$ the reconstruction efficiency, $\varepsilon_{\text{ID}|\text{Reco}}^e$ the identification efficiency for reconstructed electrons, $\varepsilon_{\text{Trigger}|\text{ID}}^e$ the trigger efficiency for identified electrons, and $\varepsilon_{\text{Iso}|\text{Trigger}}^e$ the efficiency of the isolation criterion applied to the electrons selected by the single electron trigger. Trigger efficiencies are measured before deriving the isolation efficiency, since isolation criteria are applied by the trigger, which could re-introduce differences between simulated and observed events after correcting for differences in isolation efficiency.

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

Scale factors for electron reconstruction and electron identification are centrally computed using in CMS the tag-and-probe method. The Z-boson peak is employed as narrow resonance, using the mass window: $60 \text{ GeV} < M(ee) < 120 \text{ GeV}$. Figure 8.7 shows the scale factors for (a) reconstruction efficiency, and for (b) identification efficiency using the triggering MVA ID and the working point wp80. The scale factors for the reconstruction efficiency are about 1. There is a small efficiency drop for the transition region. For electrons with $p_T > 20 \text{ GeV}$, the identification scale-factors are in the range of $0.98 - 1$, except in the endcap $|\eta| \gtrsim 1.5$, where they are in the range of $0.90 - 0.97$.

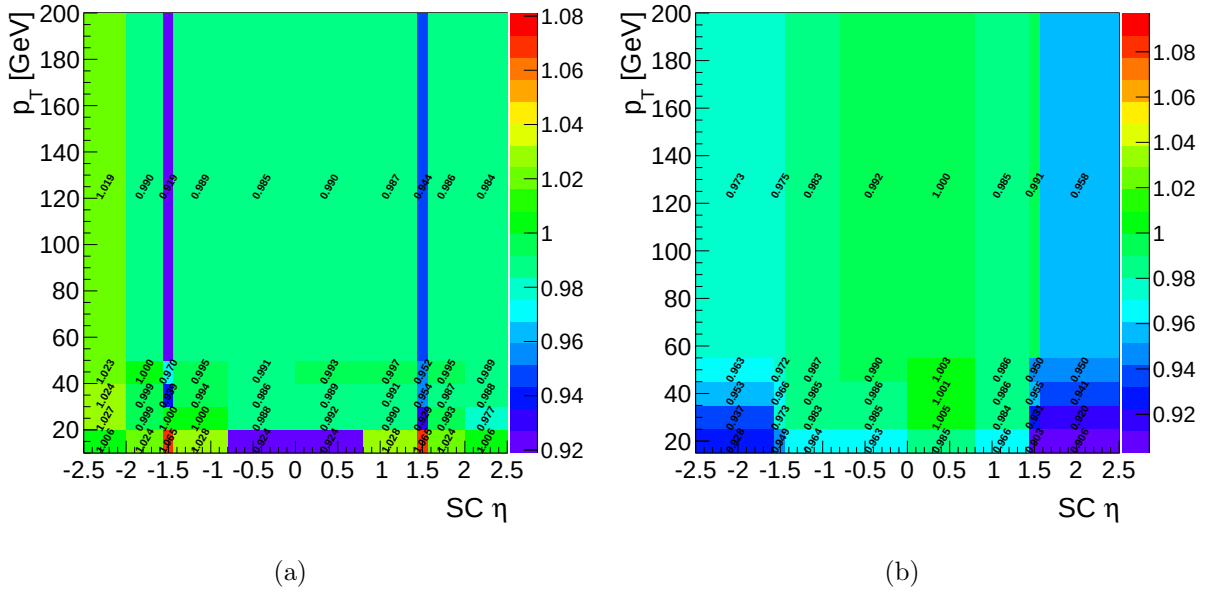


Figure 8.7.: Electron scale-factors for (a) reconstruction, and (b) identification efficiency (triggering MVA ID wp 80) as function of electron SC pseudorapidity η , and electron p_T .

Trigger and isolation scale-factors are computed using the same tag-and-probe method. Events have been recorded with the single electron trigger described in Section 8.4.1. Simulated $DY \rightarrow ll$ events were generated with MadGraph, for the matrix elements, and PYTHIA 8, for hadronisation. Efficiencies are measured in bins of SC pseudorapidity η and electron momentum. A different binning is used for the trigger and isolation efficiency-measurement to have more statistics in each bin. The same η -binning is used, but bins with the same absolute value are merged in one bin. For the reconstruction and identification efficiency-measurements, different p_T -binnings were used. The binning of the reconstruction measurement with a few new bins is used, since the other binning has low statistics in the first bin.

First, scale factors for the single electron trigger are derived. All criteria of the previous steps are applied. Tag-electrons are required to have fired the single electron trigger. Trigger efficiencies are measured using the probe electron. Scale factors are derived by

comparing efficiencies in collision data with the ones from simulation. Then, isolation criteria are applied to the tag electron and isolation efficiencies with corresponding scale factors are derived. Figure 8.8 shows the scale factors for (a) the trigger efficiency of the single electron trigger, and the (b) isolation efficiency for applying tight isolation criteria. Trigger scale-factors are between 0.77 and 1. However, the scale factors are below 0.94 only for $p_T < 30$ GeV. The single electron trigger is not fully efficient in this transverse-momentum region. Differences between efficiencies from collision data and simulated ones are largest in this region. The isolation scale factors are about 1.

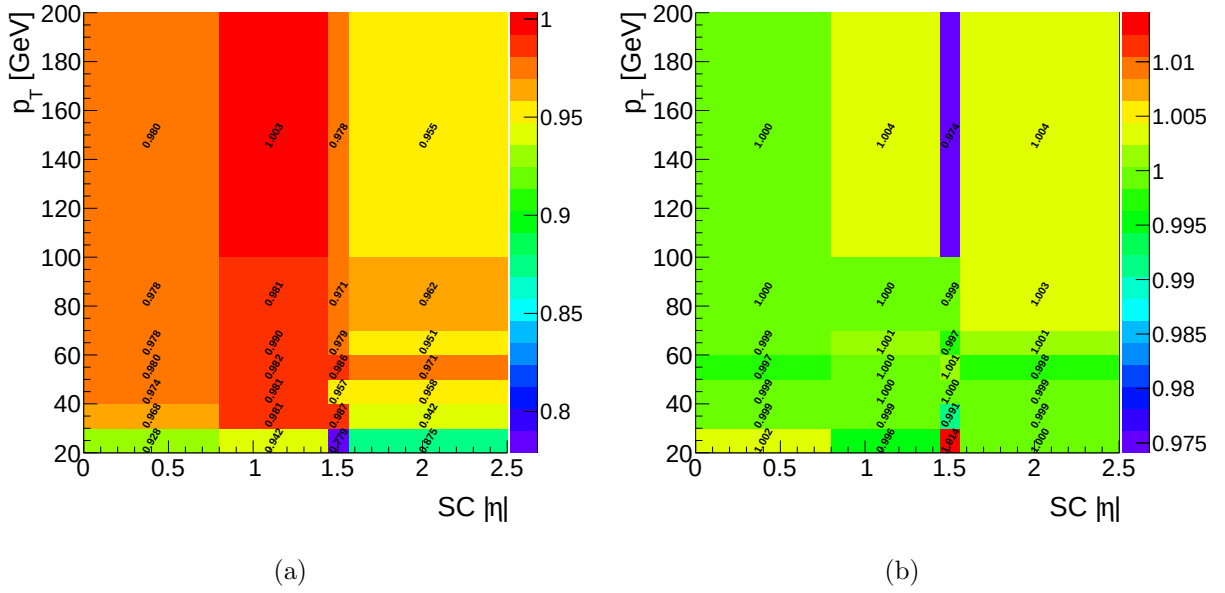


Figure 8.8.: Electron scale-factors for (a) the HLT_Ele23_WPLoose_Gsf trigger, and (b) the tight isolation as function of electron SC pseudorapidity $|\eta|$, and electron p_T .

Muon scale factors Scale factors for muon identification, muon isolation criteria, and muons triggers are centrally computed in CMS using a tag-and-probe method. The total muon identification efficiency ε^μ can be factorised as following:

$$\varepsilon^\mu = \varepsilon_{\text{ID}|\text{TRK}}^\mu \cdot \varepsilon_{\text{Iso}|\text{ID}}^\mu \cdot \varepsilon_{\text{Trigger}|\text{Iso}}^\mu, \quad (8.5)$$

where $\varepsilon_{\text{ID}|\text{TRK}}^\mu$ is the muon-identification efficiency for tracks, $\varepsilon_{\text{Iso}|\text{ID}}^\mu$ the isolation efficiency for identified muons, and $\varepsilon_{\text{Trigger}|\text{Iso}}^\mu$ the trigger efficiency for muons passing the isolation requirement. Since no muon triggers are used, the last efficiencies and corresponding scale factors are not necessary.

Studies from the Tracking POG showed that tracking muon-efficiency scale-factors for the 2015 collision data are compatible with 1. The other scale factors are derived with a tag-and-probe technique, which employs the Z boson peak as narrow resonance. A mass window

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

of $70 \text{ GeV} < M(\mu\mu) < 130 \text{ GeV}$ was used for the identification-efficiency measurement, and a mass window of $77 \text{ GeV} < M(\mu\mu) < 130 \text{ GeV}$ was used for the isolation-efficiency measurement. Events have been recorded with a single muon trigger, IsoMu20, which has requirements on the muon isolation and a requirement on the muon momentum, $p_T > 20 \text{ GeV}$. Generated events also had to pass the trigger requirements of the single muon trigger. Tag-muons were required to pass the single muon trigger requirements, to have a muon momentum larger than 21 GeV, and to have a tracker-plus-calorimeter combined relative isolation of $I_{comb}^{rel} < 0.2$, for a cone with $\Delta R < 0.4$ and after applying $\Delta\beta$ pileup-corrections. Identification efficiencies were measured using general tracks with a $p_T > 20 \text{ GeV}$ as probe muons. Isolation efficiencies were derived by using muons which pass the medium muon ID and have a transverse momentum larger than 20 GeV. Figure 8.9 shows the scale factors as function of muon p_T and $|\eta|$ for (a) identification efficiencies of the medium muon ID, and for (b) tight-isolation efficiency. They were centrally computed. Identification-efficiency scale factors are between 0.96 and 1, while scale factors for isolation efficiency are compatible with 1. There are no scale factors for muon transverse momenta smaller than 20 GeV. For isolation efficiency, corrections are very small, and thus can be neglected. There is a small dependency of the identification-efficiency scale-factor on muon p_T , therefore the corresponding scale factor for muons which have a transverse momentum of 20 GeV is used for muons with smaller transverse momenta. However, this dependency is at the per-cent level and smaller than the η -dependency,

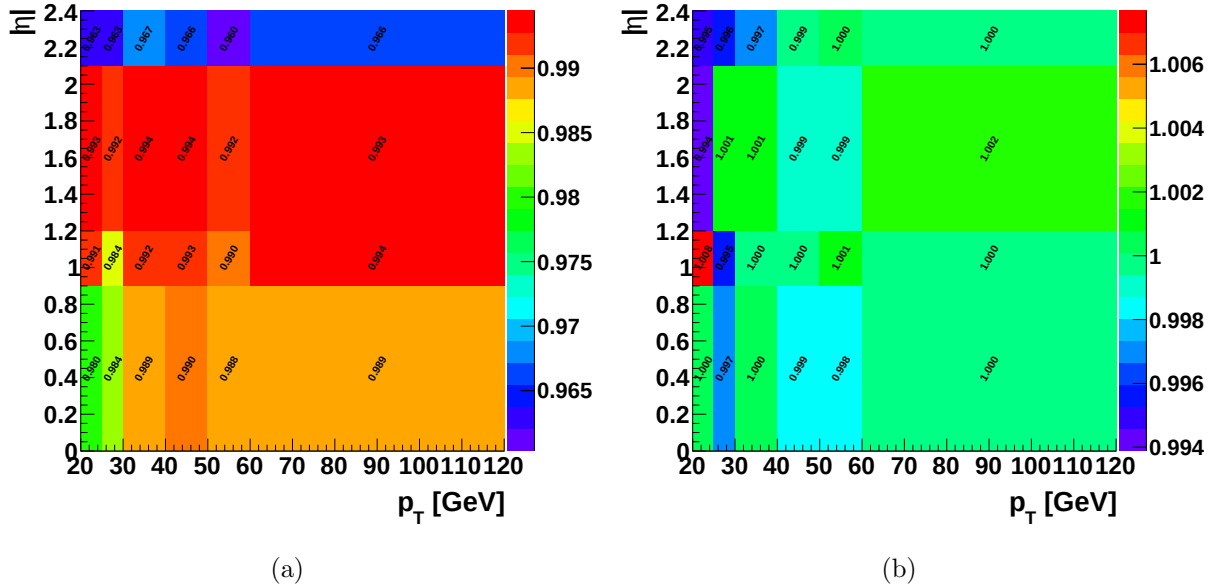


Figure 8.9.: Muon scale-factors for (a) identification (medium ID), and (b) tight-isolation efficiency as function of muon p_T and $|\eta|$. These scale factors were centrally computed using a tag-and-probe method.

Jet energy corrections Jet energy corrections, as described in Section 6.6.3, are applied to PF jets used in the analysis. Pileup offset energy corrections are applied. The jets are corrected to the simulated particle response, and residual data/simulation scale factors account for any residual differences between data and simulation after applying the other corrections including pileup reweighting.

No corrections for jet energy resolution are applied, because JER is negligible for this analysis. Only leptonic decays are investigated in the signature and jets are just used to categorise events with different production modes and background sources.

Type-I corrections for missing transverse energy Type-I corrections, described in Section 6.7.2 are applied on $\vec{E}_T$ to propagate the jet energy corrections to the $\cancel{E}_T$ reconstruction. Jets are corrected if they have less than 0.9 of their energy in the ECAL.

Scale factor for identification of b-quark jets Scale factors for the identification efficiency of b-quark jets and misidentification efficiencies for light-flavour jets, as described in Section 6.8.3, are centrally computed in CMS for jets identified as b-quark jets. The signature of LFV decays of the Higgs boson to $e\tau_\mu$ has no b-quark jets. There are no centrally computed scale factors for vetoes on jets identified as b-quark jets.

In LHC Run II, a new reweighting technique has been developed which calibrates the full discriminator shape, Section 6.8.4, and scale factors were derived using $t\bar{t}$ events. These scale factors are easier to apply on analysis which use a veto on b-quark jets. This method has been tested for this analysis, but it does not correct differences in the discriminator shape of the CSVv2 algorithm in the kinematic region of this search.

Differences in a veto of jets identified as b-quark jets between data and simulation result in normalisation differences between simulated backgrounds and observed ones. This will later be covered by the systematic uncertainties used for the normalisation of simulated events.

8.4.4. Preselection

The preselection starts by doing an event cleaning. First a matching is done between electrons and muons. If there is a muon within a cone of $\Delta R < 0.1$, the electron is removed. Jet cleaning is done after the lepton cleaning. Jets are matched to electrons and muons using a cone of $\Delta R = 0.4$, which is the jet radius parameter. If a match with a lepton is found, the jet is removed from the event.

Events are required to have an isolated electron and an isolated muon of opposite charge. Only events which have no further leptons with transverse momenta larger than 10 GeV are considered. The single electron trigger is also applied on simulated events and the electron is required to fire the trigger. A matching between the reconstructed and the trigger electron object is done to check if the electron has fired the trigger. The electron is required to have $p_T^e > 30$ GeV, which is a tighter requirement on the transverse momentum than the electron object definition, $p_T^e > 20$ GeV. This tighter criterion was chosen, since the single electron trigger is not fully efficient for lower transverse momenta and differences between efficiencies from collision data and simulated ones are the largest for $p_T^e < 30$ GeV. Furthermore, the amount of misidentified electrons is suppressed.

As the signatures of LFV Higgs-boson decays do not include b-quarks, a veto on events with jets identified as b-quark jets is applied to reduce the overwhelming background of events with b-quark jets originating from top-quark decays. The visible mass, defined as the invariant mass of the electron and muon, has to be larger than 12 GeV to suppress backgrounds from low mass resonances. Only events with less than three jets are considered.

Events are classified into three exclusive jet categories, which are sensitive to the different Higgs-boson production-mechanisms. The 0-jet category and the 1-jet category are sensitive to GF Higgs boson production, while the 2-jet category is sensitive to both, GF and VBF Higgs-boson production.

8.5. Background modelling

All contributions from background processes are estimated using simulations. The normalisation of the individual background contributions is checked and systematic uncertainties are derived in control regions, described in Section 8.6. Then, the background model is fitted to collision data within the systematic uncertainties of the model. After applying the preselection requirements, the dominant background processes are the SM $Z \rightarrow \tau\tau$ production and misidentified leptons, which predominantly are found in W+jets events. QCD multijet production contributes to the misidentified-lepton background, but contributions are small because of the small probability that two PF objects are misidentified, one as an isolated electron and the other one as an isolated muon with opposite sign. Another important type of background processes, for events with one or two jets, is the top

background, where the $t\bar{t}$ and tW processes contribute to. The contributions from both processes are reduced by the veto on jets identified as b-quark jets, but not all b-quark jets are identified by the CSVv2 algorithm. The diboson processes, WW , WZ , and ZZ +jets, are combined to another background. Background contributions from both SM Higgs boson production processes, described before, the $Z \rightarrow ee/\mu\mu$ processes, and the $W\gamma/\gamma^*$ processes are also considered. SM $H \rightarrow \tau_e\tau_\mu$ processes only contribute to a small amount. Their collinear mass distribution is very similar to the signal one, therefore it has to be taken care, that such kind of events are negligible. As both leptons originate from tau decays, the electron momentum spectrum is softer than for LFV decays of the Higgs boson to $e\tau_\mu$. Additionally, there are only neutrinos from one tau decay, which are collinear to the muon flight-direction. Thus, the SM Higgs-boson background can be distinguished from LFV Higgs-boson decays and it can be reduced to a negligible amount.

Used simulated samples and their cross sections are listed in Table B.1 and Table B.2 from Appendix B, respectively. The contributions from the dominant background processes, SM $Z \rightarrow \tau\tau$ production and misidentified leptons, can be estimated using collision data. This was done for the $\sqrt{s} = 8$ TeV analysis [5]. However, simulation only is used for this analysis. The technique for estimating the $Z \rightarrow \tau\tau$ background cannot be used for this analysis because a corresponding sample was not centrally produced by CMS for the 2015 collision data. As the statistic of the 2015 collision data-sample was lower than for the $\sqrt{s} = 8$ TeV data sample, estimating the misidentified-lepton background from collision data describes the collision data in a control region as well as the MC simulation, this will be described later. Thus, simulation was used to model this background and the event selection will be optimised to suppress this background to a negligible amount. Nevertheless, both methods are described in the following.

8.5.1. $Z \rightarrow \tau\tau$ background

The $Z \rightarrow \tau\tau$ background contributions can be estimated using an embedding technique [126] which employs lepton universality. Key features of the event topology are taken from collision data, such as the jet multiplicity N_{Jet} , $\vec{E}_T$ and underlying event, while the simulation is used for modelling the detector response of the tau decay-products.

For each data-taking period, $Z \rightarrow \mu\mu$ events, which have a very clean signature, are recorded applying a loose $Z \rightarrow \mu\mu$ selection. Then, Muons are replaced with simulated tau decays. This is done by replacing the PF muons by the PF particles which are reconstructed from visible decay products of the taus in simulated $Z \rightarrow \tau\tau$ events. Afterwards, jets, hadronic taus, and $\vec{E}_T$ are reconstructed and the lepton isolation is calculated.

Scale factors are derived to correct the identification and isolation efficiencies of the embedded sample to the one expected from simulation. Furthermore, there is a bias in the embedded sample relative to simulation due to the effect of final-state radiation of

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

photons. This leads to a shift of the mass peak in the visible mass distribution. In the analysis using the $\sqrt{s} = 8$ TeV collision data, this shift was found to be in the order of 2% [5]. The normalisation of the embedded sample is taken from simulation by comparing the event yield in the visible mass distribution. These corrections are derived with the visible mass, since electrons and muons have clean signatures in the detector.

The embedding technique is not used in this analysis because the embedded sample was not available for the 2015 data. Thus, the $Z \rightarrow \tau\tau$ background process is taken from MC simulation as the other $Z \rightarrow ll$ backgrounds.

8.5.2. Misidentified lepton background

The misidentified lepton background is dominated by the W+jets process, where the W boson decays to a lepton and a neutrino, and another PF candidate is misidentified as the second lepton, denominated *fake lepton*. Other processes which contribute to the misidentified lepton background are discussed in Section 8.3. The misidentified background contribution has the highest systematic uncertainty, which is in the order of 40-50%¹. Isolation requirements can reduce the amount of misidentified leptons, since such leptons mostly arise from jets. Figure 8.10 shows the composition of the fake objects from W+jets simulation as function of (a) p_T^e , (b) p_T^μ , and (c) N_{Jet} normalised to the expected number of events after preselection for an integrated luminosity of 2.3 fb^{-1} . The relative fraction of muon fakes and electron fakes depends on the electron and muon p_T . There is also a small dependence on the jet multiplicity. As a higher p_T requirement is applied on the electron than on the muon, the misidentified lepton background is dominated by PF candidates which are misidentified as muons. Thus, the amount of misidentified electrons can be further reduced by applying a tighter criteria on the electron transverse momentum, as shown in Figure C.1 from Appendix C. This will be employed later in the optimisation of the event selection.

¹The systematic uncertainty on the misidentified lepton background is estimated in a control region described in Section 8.6. In the search at $\sqrt{s} = 8$ TeV and in this search, the systematic uncertainty was found to be 40%. It is also the same for the background estimation using collision data and for the estimation using MC simulation.

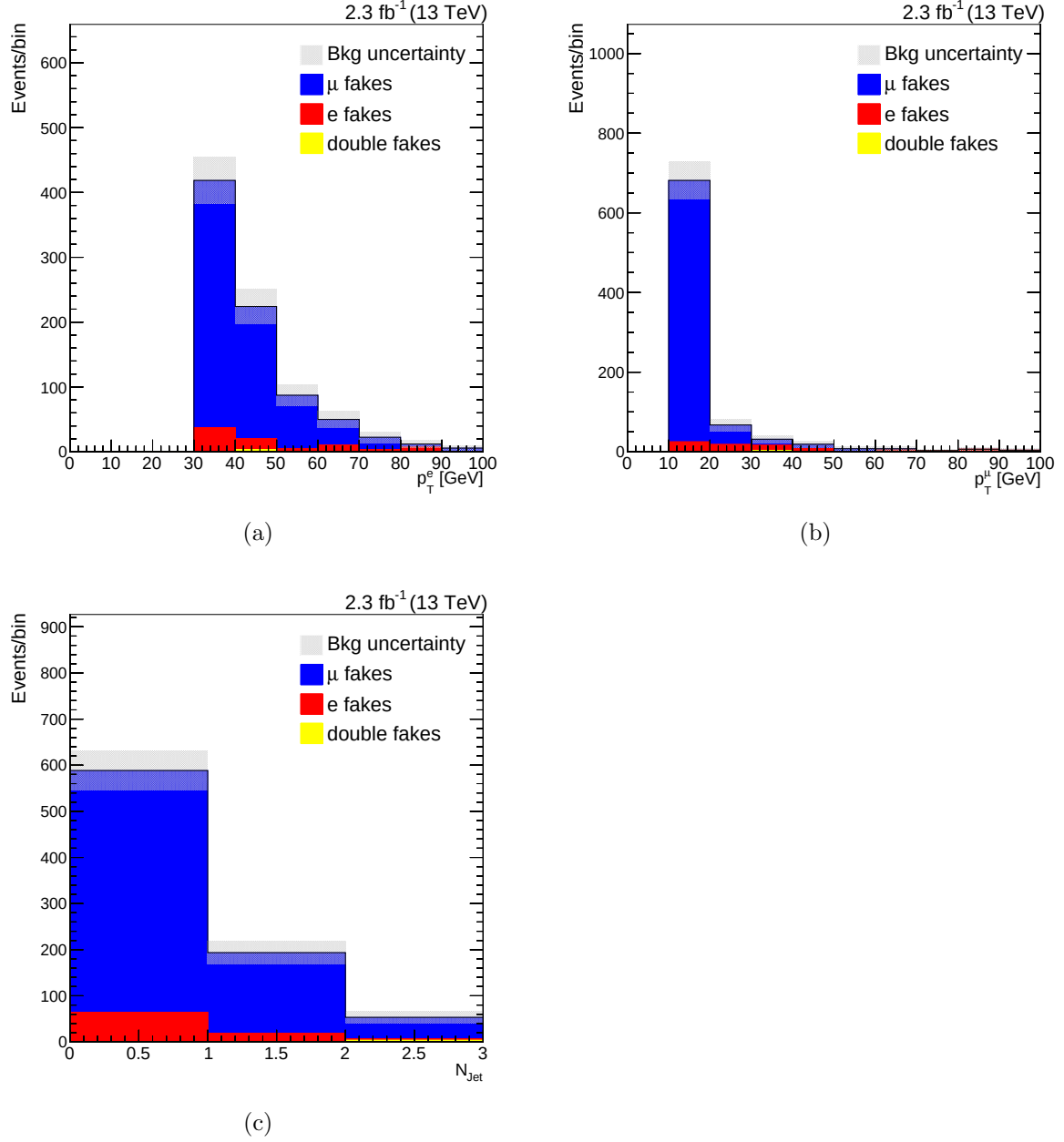


Figure 8.10.: Composition of the misidentified lepton background after preselection from W+jets simulation as function of (a) electron p_T , (b) muon p_T , and (c) jet multiplicity. The background uncertainty corresponds to the statistical uncertainty of the MC simulation.

Isolation side-band method

A tighter requirement on the isolation reduces the amount of misidentified leptons. Furthermore, inverting the isolation requirement on one of the leptons defines a region enriched with misidentified leptons. As the amount of misidentified electrons and double misidentified leptons is negligible in comparison to misidentified muons, the relative muon PF isolation is inverted. Isolated muons have a relative PF isolation smaller than 0.15, while non-isolated muons have a relative PF isolation of $0.15 < I_{\text{PF}}^{\text{rel}} < 1$. This isolation side-band includes muons with looser isolation values than the loose-isolation criterion of $I_{\text{PF}}^{\text{rel}} < 0.25$. Muon-isolation scale-factors were centrally computed only for muon isolation values tighter than the loose-isolation. Scale factors for the tight and the loose isolation are in the order of 1, as one can see in Figure 8.9 and Figure C.2 from Appendix C, therefore the effect of using no scale factors for $0.25 < I_{\text{PF}}^{\text{rel}}$ is negligible.

Another misidentified-lepton enriched region can be defined by requiring same charge of the leptons because the SM rate for true same sign leptons is very small. Altogether, four regions can be defined, as shown in Figure 8.11, where region I is the signal region with

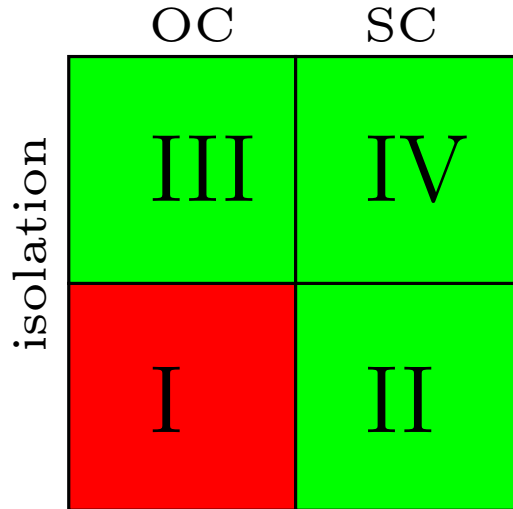


Figure 8.11.: Misidentified lepton regions

two isolated leptons of opposite charge. In region II, the charge requirement is inverted, so the leptons have the same charge. This region can be used to validate the method used for estimating the misidentified-lepton background, Figure C.3 from Appendix C. The other two regions have inverted muon-isolation requirements. Region III has two leptons with opposite charge, an isolated electron and a non-isolated muon, while region IV has an inverted requirement on the lepton charges with respect to region III. The ratio of isolated

to non-isolated muons depends on muon p_T and η . It is defined as:

$$f_\mu(p_T, \eta) = \frac{N_\mu^{\text{iso}}}{N_\mu^{\text{non-iso}}}[p_T, \eta], \quad (8.6)$$

where N_μ^{iso} is the number of isolated muons, and $N_\mu^{\text{non-iso}}$ is the number of non-isolated muons. When f_μ is known, the event yield of misidentified muons can be estimated in region I by taking the event yield in region III and reweighting it with f_μ . Thus, the misidentified-lepton background can be estimated from collision events. There are other background processes which include prompt non-isolated leptons, or misidentified leptons, such as $t\bar{t}$ pair-production and diboson processes. They are already taken into account in other background sources and must not be double counted. Contributions of all other background processes to this method are subtracted from the event yield of misidentified leptons using MC simulation. The other two regions with same charge leptons can be used as cross check.

The misidentification ratio f_μ can be measured in $(Z \rightarrow \mu\mu) + X$ events, where X is either an isolated or a non-isolated muon. Contributions from SM background processes are small. Diboson, $t\bar{t}$, tW , and the $W\gamma/\gamma^*$ processes contribute to this background and they are subtracted. Collision events were collected using double-muon triggers, while the event yield of the other background processes is taken from simulation.

Events which have three muons and no further leptons are selected. Only events with less than three jets and none of them identified as a b-quark jet are further considered. It is required that there are two isolated muons with opposite charge, which have a transverse momentum larger than 25 GeV. The invariant mass of both muons has to be close to the Z-boson mass, $80 \text{ GeV} < M(\mu\mu) < 100 \text{ GeV}$. Finally, a requirement on the missing transverse energy of $\cancel{E}_T < 50 \text{ GeV}$ is applied to select mostly events with a Z boson decaying to two muons. This criterion is necessary to avoid for some kinematic regions in which the background event-yields are larger than those for $(Z \rightarrow \mu\mu) + X$ events. The missing transverse energy in W+jet events is expected to be dominated by the neutrino energy. Replacing the W boson with a Z boson shifts the $\cancel{E}_T$ distribution to smaller values. Thus, there should be only a very small bias by applying a requirement on $\cancel{E}_T$, as it can be seen from Figure C.4 from Appendix C.

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

Figure 8.12 shows the measured misidentification ratio, which can be measured with a relative statistical uncertainty of about 15% for muons with $10 \text{ GeV} < p_T < 20 \text{ GeV}$. For muons with $20 \text{ GeV} < p_T < 40 \text{ GeV}$, the relative statistical uncertainty is in the order of 40% in the barrel and $\sim 75\%$ in the endcap. All other bins have relative uncertainties in the order of $\sim 80\%$. Misidentification ratio are between 0.36 and 0.4 in the barrel region, and in the order of 0.55 in the endcap region, for bins where the relative uncertainties are below 50%. As most misidentified muons have a p_T smaller than 20 GeV, as Figure 8.10(b) shows, high uncertainties in the misidentification ratios for other transverse momentum regions have a negligible contribution to the overall uncertainty of the misidentified-lepton background.

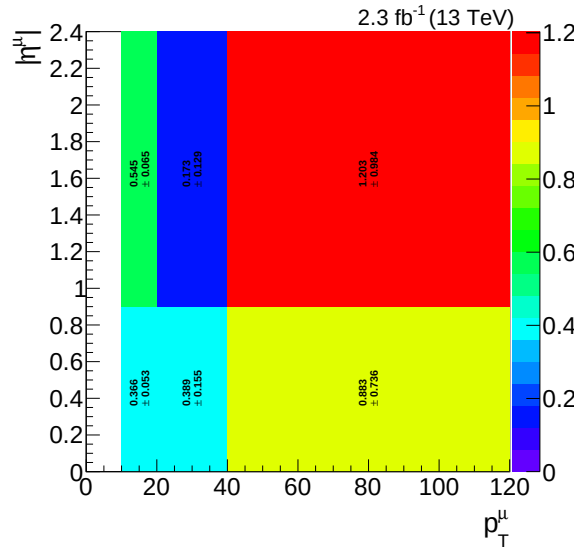


Figure 8.12.: Misidentification ratio f_μ as function of muon p_T and η . The misidentification ratio was derived in $(Z \rightarrow \mu\mu) + X$ events, where X is either an isolated or a non-isolated muon.

Validation The isolation side-band method and the misidentification ratios can be validated using the same-charge regions II and IV, shown in Figure 8.11. Region IV is dominated by misidentified leptons. Applying the misidentification ratio in this region gives an estimate for the misidentified-lepton background contribution in region II. Several validation distributions are shown in Figure C.5 of Appendix C. Figure 8.13 shows the collinear mass after preselection using MC simulation (left) or the isolation side-band method (right) for estimating the misidentified-lepton background. The misidentified-lepton background is slightly underestimated using MC simulation, while it is overestimated using the isolation side-band method. This is covered by the systematic uncertainty of the misidentified-lepton background.

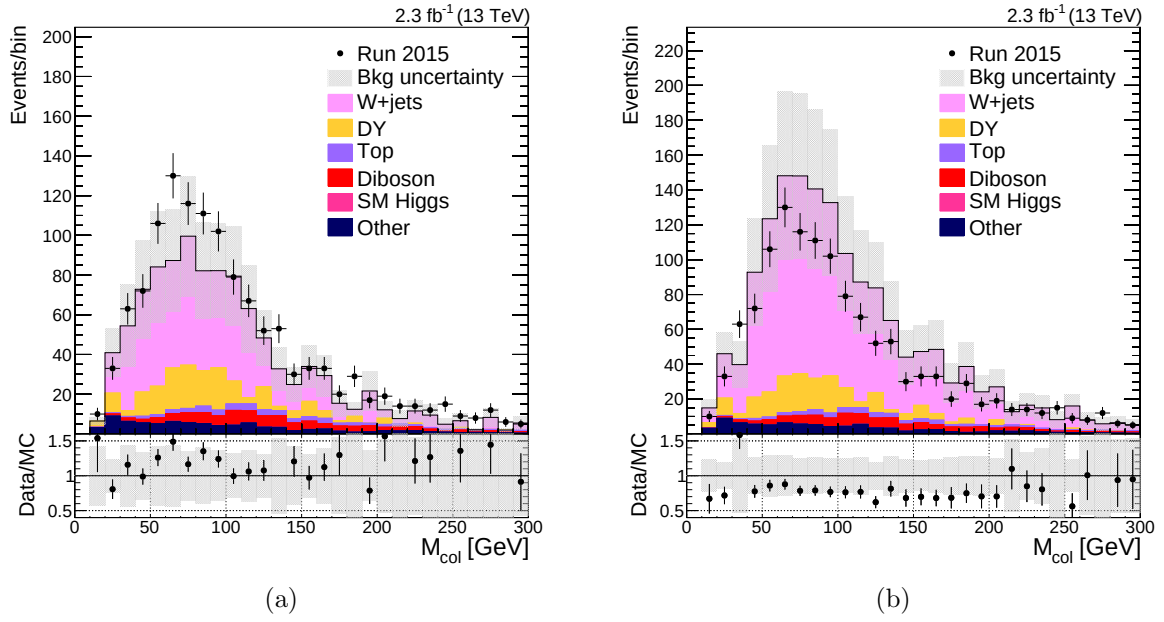


Figure 8.13.: Collinear mass after preselection in the same-charge region. The misidentified-lepton background was estimated using (a) MC simulation, or (b) the isolation side-band method.

Methods to estimate the misidentified-lepton contribution

The misidentified-lepton background can be estimated using different methods. In the following, assumptions and drawbacks are discussed. Many events have to be generated to have enough statistic for estimating the misidentified-lepton background contribution with MC simulation. These events have usually large weights. Simulated misidentified-leptons have a high systematic uncertainty, therefore such background processes are usually estimated using collision data, when the contribution of this background is large and cannot be neglected. The systematic uncertainty is in the order of 40% with an extra uncertainty of 10% to account for normalisation differences dependent on the jet multiplicity.

Isolation side-band method using $(Z \rightarrow \mu\mu) + X$ events: This method employs the isolation side-band for estimating the misidentified-lepton background and has already been discussed. The misidentification ratio f_μ is measured in $(Z \rightarrow \mu\mu) + X$ events, collected with double-muon triggers. It assumes that the W boson can be replaced by a Z boson. Since the probability that a PF candidate is misidentified as a lepton should be independent of the presence of the additional vector boson, and most differences cancel each other when deriving the misidentification ratio, this assumptions holds.

There could be a small bias due to the muon triggers, since it is possible that the third muon and only one muon from the Z-boson decay, instead of both, fired the trigger. In this case, differences in trigger efficiency would not necessarily cancel, as the trigger efficiency depends on the muon p_T , η , and isolation. This would affect mainly muons with higher p_T , or smaller values for isolation, where the statistics of the 2015 collision data is low.

As other background contributions have to be subtracted when measuring the misidentification ratio f_μ , this background has to be reduced by selection criteria. No requirements should be used which select kinematic regions that differ from the regions used as signal region and isolation side-band region. The isolation side-band method also assumes that the composition of misidentified leptons is the same in both isolation regions. This does not have to be the case, therefore a systematic uncertainty on the misidentified-lepton background of 40% and an extra uncertainty for each jet category of 10% to account for differences in the normalisation in dependence on the jet multiplicity are used.

Isolation side-band method using $(Z \rightarrow ee) + X$ events: This method employs $(Z \rightarrow ee) + X$ events to measure the misidentification ratio f_μ . Collision data collected with single electron triggers or double-electron triggers can be used. Using the same trigger as for the signal region, events which are very similar as the signal region will be selected, except for the second electron, and any trigger biases on the muon are removed. This was tested, but the statistic was too low for measuring the misidentification ratios using the 2015 collision data. A data sample with more statistic is needed to check if using a sample which was recorded using the same trigger thresholds can reduce the systematic uncertainty for the misidentified-lepton process. Assumptions and systematic uncertainties are very similar to the previously described method for measuring the misidentification ratio f_μ .

All discussed methods for estimating the misidentified-lepton background, including the estimation from MC simulation only, have high systematic uncertainties. Using collision data for estimating the background would remove the systematic uncertainty from simulation. Statistical uncertainties would also become smaller. On the other hand, systematic uncertainties originating from knowledge of the composition and origin of misidentified leptons have to be reduced. Thus, it was studied if the composition of misidentified leptons is similar in all four regions. Furthermore, properties of the misidentified-lepton background were studied in each region and compared to each other.

Composition of the misidentified-lepton background in each region Figure 8.14 shows the composition of the misidentified-lepton background in MC simulation of W+jets events after preselection for each region as function of the collinear mass. Similar distributions are shown for the jet multiplicity in Figure C.6 of Appendix C. Most misidentified leptons are misidentified muons. There is a small contribution of misidentified electrons for regions I, II with two isolated leptons, which in size is comparable to the statistical uncertainty. In regions with a non-isolated muon (regions III, IV), there is a negligible contribution from misidentified electrons. Thus, the composition of misidentified leptons differs slightly in the isolation side-bands from the regions with two isolated leptons. The misidentification ratio covers only the misidentified muons. It is more difficult to estimate the misidentified electron background using the isolation side-band method, since an isolation requirement is applied as part of the electron trigger, which is in the order of $I_{\text{PF}}^{\text{rel}} < 0.2$. Either the misidentification ratio has to be measured applying this trigger or the misidentification background has to be estimated using a collision data sample which was collected without trigger requirements on the electron. The first is possible, but would further reduce the statistics, the second method is not possible due to the small transverse momentum of the muon, for which a single muon trigger does not exist. It was checked if an isolation side-band $0.1 < I_{\text{PF}}^{\text{rel}} < 0.2$ has enough statistic for $\sqrt{s} = 8$ TeV collision data and for $\sqrt{s} = 13$ TeV collision data. The statistics in both samples is too small for this isolation side-band. Uncertainties would be even higher than neglecting this contribution.

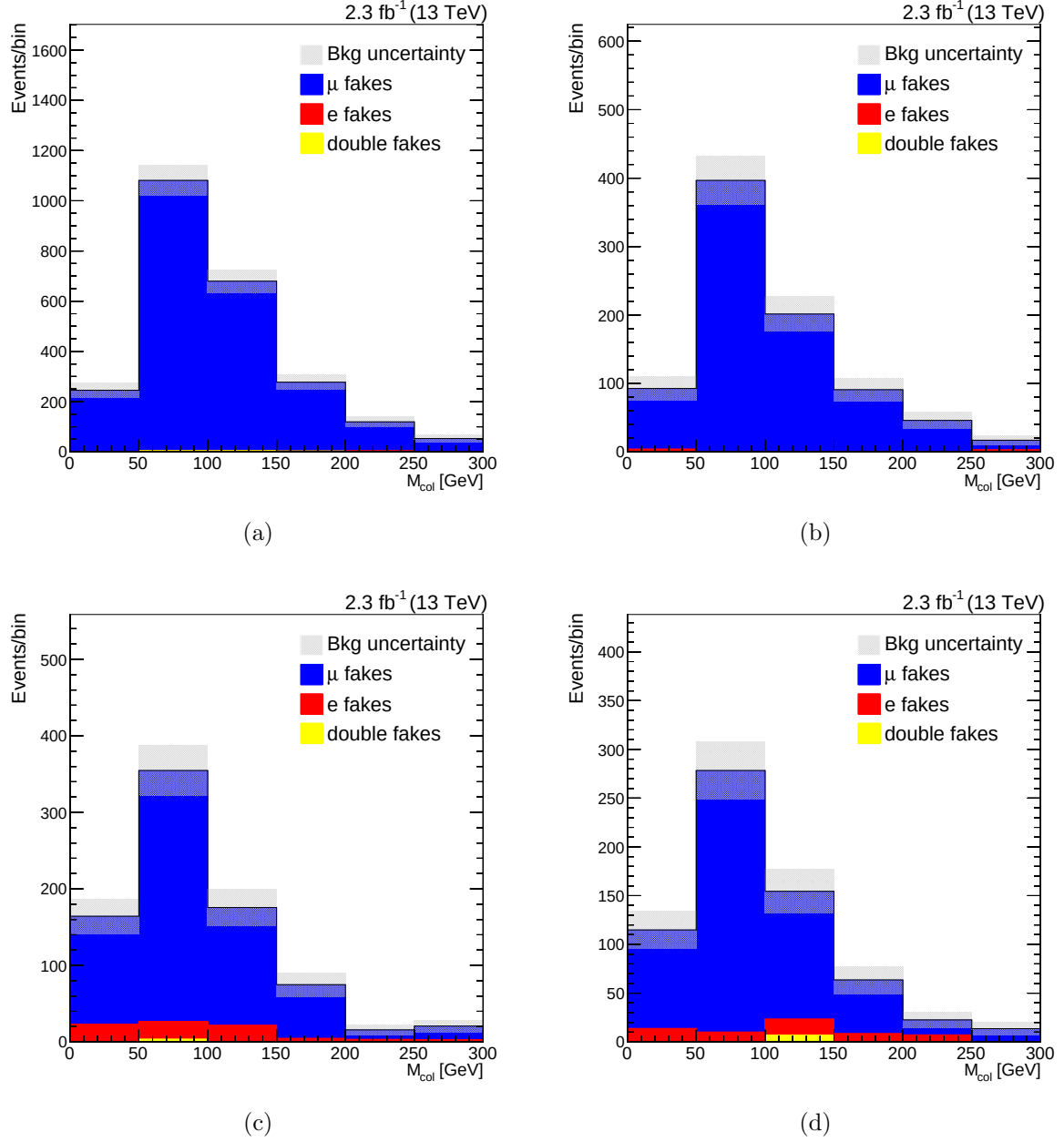


Figure 8.14.: Composition of the misidentified-lepton background as function of the collinear mass after preselection in (a) region III, (b) region IV, (c) region I, and (d) region II. Events were generated using MC simulation.

Sources for misidentified leptons Misidentified muons can either originate from real muons, which are non-prompt, or from other PF candidates. Figure 8.15(a) shows the muon sources for each of the four categories. Only simulated events passing the preselection and where the W boson decays to an electron and corresponding neutrino are considered. Both regions with a non-isolated muon (region III and region IV) have a higher fraction of events with real muons than the regions with two isolated leptons (region I and region II). The misidentification ratio for opposite-charge and same-charge leptons is shown in Figure 8.15(b). The misidentification ratio is much higher for fake muons than for real muons. Furthermore, there is a slight dependence on the charges of both leptons.

As misidentification ratio depends on the process which is not known in collision data, it has to be measured in an independent sample of collision data which has similar composition of these fake muon sources, or one of these sources has to be suppressed by selection criteria. Finding an independent sample which has the same composition of muon sources as the signal region is difficult, since there are dependencies on other observables, such as if both leptons have opposite or same charges. Figure C.7 from Appendix C shows dependencies on further observables, which are the jet multiplicity, the collinear mass, and the angular difference in ϕ between the electron and the muon, $\Delta\phi(e, \mu)$. The dependence of the misidentification ratio on $\Delta\phi(e, \mu)$ makes it even more difficult to estimate the misidentified-muon background in the signal region using another independent region or data sample. The misidentification ratio decreases with $\Delta\phi(e, \mu)$, therefore requiring large $\Delta\phi(e, \mu)$ values can reduce the background of misidentified muons.

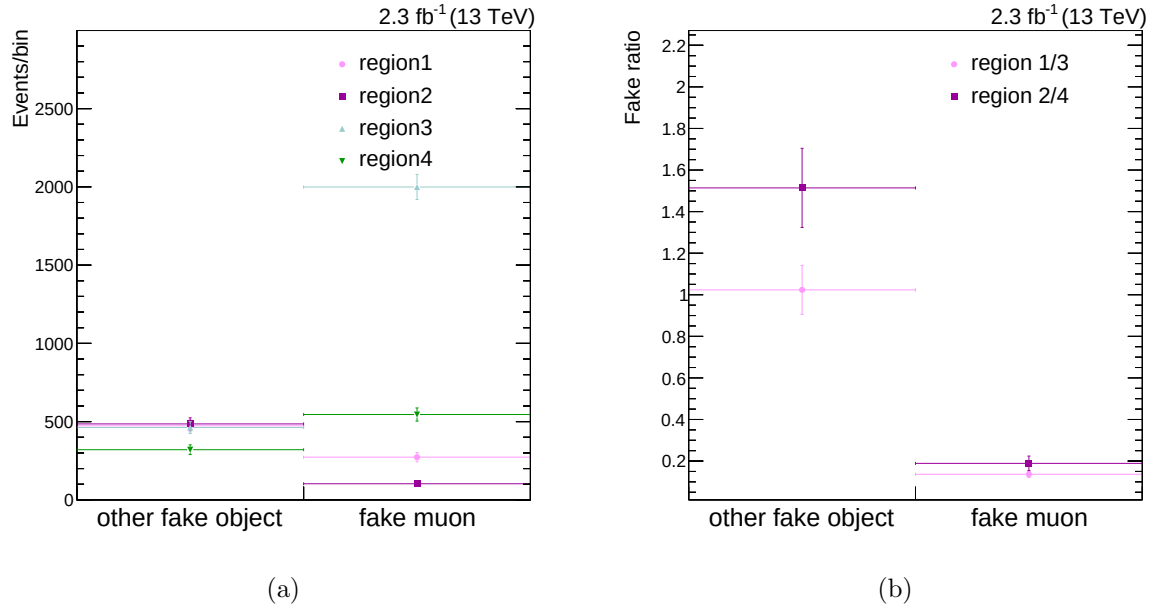


Figure 8.15.: Sources for misidentified muons are compared between the four regions using MC simulation. Events are classified according to the selected muon. Events with real muons are combined to the fake muon category, while contributions due to other PF candidates are classified as fake muon. The number of events for each category in each region is shown in (a) and (b) shows the ratio of isolated to non-isolated muons as function of each category for the opposite-charge region (I/III) and the same-charge region (II/IV).

Conclusions

Several methods were described for estimating the misidentified lepton background. Samples from MC simulation suffer from low statistic and high systematic uncertainties. On the other hand, the isolation side-band method assumes that the misidentification ratio of isolated to non-isolated muons can be measured in an independent sample and that the composition of sources for misidentified leptons is very similar. This is not exactly the case. Systematic uncertainties have to account for such differences. Both methods describe the collision data within a systematic uncertainty of 40%, as shown in Figure 8.13. This systematic uncertainty and how it is derived will be discussed later in Section 8.6. Using MC simulation is technically the simplest way of estimating the misidentified lepton background and it also includes misidentified electrons. Thus, misidentified leptons are estimated using MC simulation. This background contribution has the highest systematic uncertainty compared to all other background processes. It is difficult to reduce this systematic uncertainty, therefore reducing this background to a negligible amount makes the analysis less sensitive to it. This is important when the limit on the branching fraction of LFV decays of the Higgs boson to $e\tau$ gets dominated by the systematic uncertainties of the analysis.

8.6. Background normalisation

All background processes are estimated using MC simulation. In the previous section, Section 8.5, the background modelling was discussed. This section focuses on the normalisation of the background contributions included in the total background model. The normalisation of the most important background contributions is checked in dedicated control regions, which should not overlap with the signal region. Collision data were collected using the same single electron trigger as for the signal region. Table 8.6 gives an overview of the signal region and the control regions. The regions differ in the lepton multiplicity, the relative charge of the electron and muon for the $e\mu$ signature, and the presence of b-quark jets. Only the WW+jet background process overlap with the signal region definition. A requirement on the angle between $\vec{E}_T$ and the muon flight-direction is introduced to suppress contributions from LFV decays of the Higgs boson to a negligible amount.

Table 8.6.: Overview of signal region and control regions. Lepton multiplicity, relative charges, and presence of b-quark jets are criteria for defining independent control regions.

Process	$e\mu$	2e	$e2\mu$	$2e2\mu$	opposite charge	b-quark jet veto
LFV $H \rightarrow e\tau_\mu$	✓	-	-	-	✓	✓
Drell-Yan	-	✓	-	-	✓	✓
Misidentified leptons	✓	-	-	-	✗	✓
$t\bar{t}$, $tW/\bar{t}W$	✓	-	-	-	✓	-
WW+jets	✓	-	-	-	✓	✓
WZ+jets	-	-	✓	-	-	✓
ZZ+jets	-	-	-	✓	-	✓

8.6.1. Drell-Yan: $Z \rightarrow ll$

The Drell-Yan background is estimated by changing the flavour of the second lepton. Events are selected according to a selection dedicated to $(Z \rightarrow ee) + \text{jets}$ events. Simulated events were generated with aMC@NLO, as shown in Table B.1 from Appendix B. One electron is required to have fired the same single electron trigger used for the signal region. Two electrons of opposite charge are required with an event veto on further isolated leptons. The electron with the larger transverse momentum is required to have $p_T^{e1} > 30$ GeV, and the second electron is required $p_T^{e2} > 20$ GeV. A veto on jets identified as b-quark jets is applied and only events with $N_{\text{Jet}} < 3$ are considered. Finally, a selection on the invariant mass of both electrons is required, $60 \text{ GeV} < M(ee) < 120 \text{ GeV}$.

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

Table B.2 from Appendix B lists the recommended production cross-section for the Drell-Yan process, which is $\sigma_{MC} = 5765.4$ pb. The production cross-section was measured in $Z \rightarrow \mu\mu$ events for the dimuon mass in the range 60 to 120 GeV. The measured total inclusive cross-section times branching fraction of was:

$$\sigma(pp \rightarrow ZX \times B(Z \rightarrow \mu\mu)) = 1870 \pm 2 \text{ (stat)} \pm 35 \text{ (syst)} \pm 51 \text{ (lumi)} \text{ pb [127].}$$

This corresponds to a production cross-section of 5610 pb, when all three flavours are considered, with a systematic uncertainty of 4.6%. The measured and predicted production cross-sections differ by 3%. The production cross-section of Table B.2 is used in the analysis with a systematic uncertainty of 10% which accounts for the difference to the measured value and its systematic uncertainty.

Figure 8.16 shows the invariant dielectron mass and the jet multiplicity for the Drell-Yan control region. The normalisation agrees with the collision data within the uncertainties. An additional uncertainty of 5% should be applied for each jet category, as the production cross-section depends on the jet multiplicity. Further control distributions are shown in Figure D.1 from Appendix D.

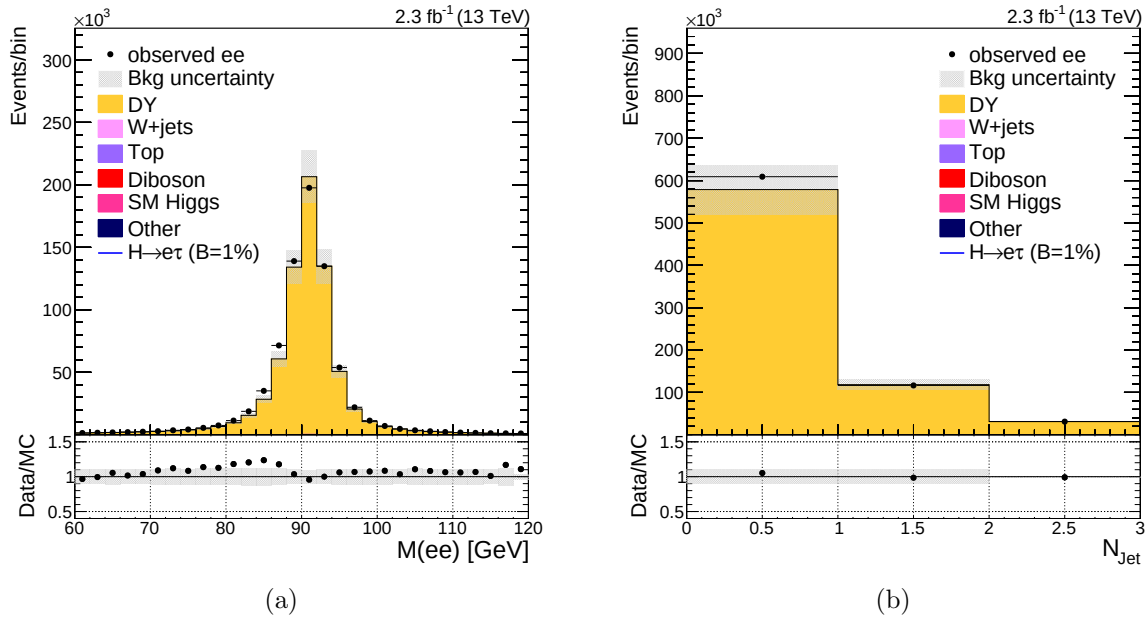


Figure 8.16.: Drell-Yan control region - The invariant dielectron mass (a), and jet multiplicity N_{Jet} (b) are shown. Collision events were collected using a single electron trigger during the 2015 LHC run. Events were generated with aMC@NLO applying a systematic uncertainty of 10%.

8.6.2. Top background

The normalisation of the top background is validated by inverting the veto on jets identified as b-quark jets. A tighter requirement on the muon, $p_T^\mu > 20$ GeV, is used to suppress background contributions from $t\bar{t}$ processes with misidentified muons. Matrix elements were generated with POWHEG and the parton shower was generated with PYTHIA 8.

In a common ATLAS-CMS effort, the top-quark-pair cross section was calculated using the Top++v2.0 program [128], which calculates the cross section in a fixed order approach with next-to-next-to leading order (NNLO) accuracy, and includes soft-gluon contributions for the hadronic cross-section with next-to-next-to-leading logarithmic (NNLL) accuracy. The predicted $t\bar{t}$ production cross-section for $\sqrt{s} = 13$ TeV and a top-quark mass of 172.5 GeV is $\sigma_{\text{MC}} = 831.76_{-29.20}^{+19.77}$ (Scale) ± 35.06 (PDF + α_S) $_{-22.45}^{+23.18}$ (Mass) pb. The systematic uncertainty is about 6%. The top-quark pair production cross-section was measured in events with one electron, one muon, and at last two jets [129]:

$$\sigma_{t\bar{t}} = 772 \pm 60 \text{ (stat)} \pm 62 \text{ (syst)} \pm 93 \text{ (lumi)} \text{ pb.}$$

The total uncertainty, without luminosity uncertainty, is 11%. The measured and estimated production cross-sections differ by $\sim 10\%$. Validation in the top-background control-regions shows, that the background estimates agree better with observed collision data when using the measured cross section, therefore it is used in further estimations of the $t\bar{t}$ process contribution. A normalisation uncertainty of 20% is used to account for the deviation from the predicted cross-section and for the total uncertainty of the measurement. This uncertainty takes also contributions of the efficiency of the identification of b-quark jets into account. Figure 8.17 shows the event yield as function of the transverse mass $m_T(\mu, \cancel{E}_T)^2$ for (a) 1 identified b-quark jet, and (b) two identified b-quark jets. In both control regions, the predicted event yields agree with the observed ones within the uncertainties. The event yield as function of the jet multiplicity in both control regions is shown in Figure D.2 from Appendix D. There is a dependence of the normalisation on the jet multiplicity. An extra uncertainty for each jet category accounts for differences between the jet categories in the event yield of the $t\bar{t}$ process.

²transverse mass $m_T(\mu, \cancel{E}_T)$ is defined in Section 8.9.2

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

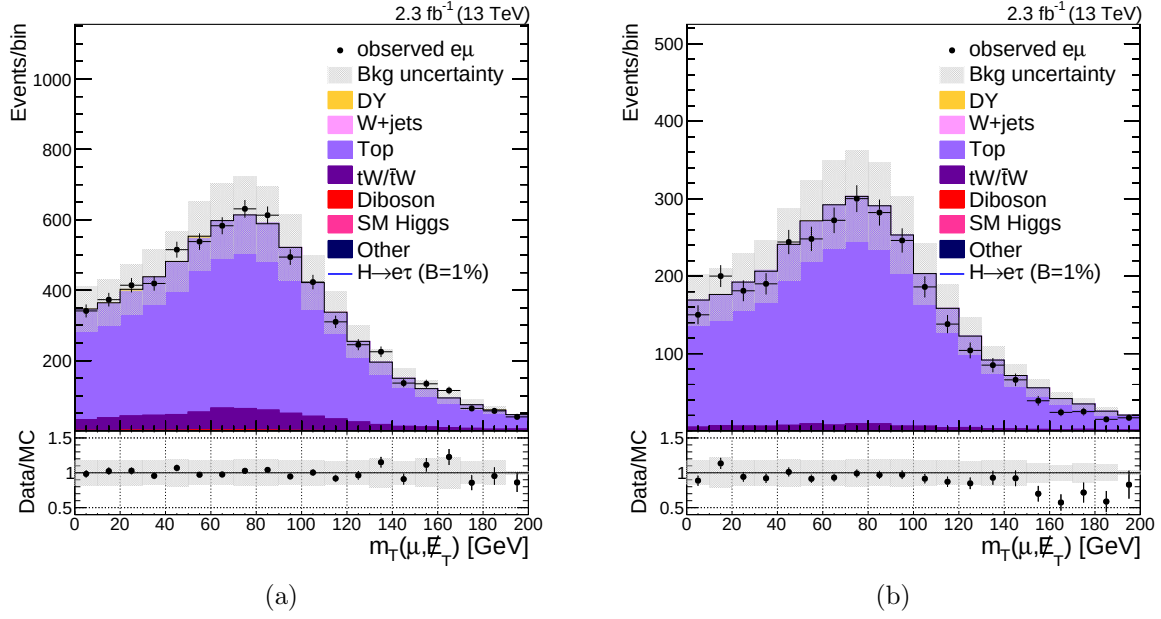


Figure 8.17.: Validation of the $t\bar{t}$ process in the top-background control-region. The event yield is shown as function of the transverse mass $m_T(\mu, \cancel{E}_T)$ for (a) 1 identified b-quark jet, and (b) two identified b-quark jets. Events were generated with POWHEG for the matrix-element generation and PYTHIA 8 for the parton shower. Simulated $t\bar{t}$ events were normalised according to the measured production cross-section [129].

8.6.3. Diboson background

The diboson contributions are from the W^+W^- -production process, the WZ -production process, and the ZZ -production process. A control region is defined for each diboson process to validate the normalisation of each individual background. Then, the composition of the diboson background contribution and its total uncertainty is discussed.

W^+W^- process

The control region for the WW +jets process requires an isolated electron and an isolated muon of opposite charge. A veto on further isolated leptons and jets identified as b-quark jets is applied. Further criteria are applied to reduce the overlap with the signal region and to suppress other background processes. Both leptons are required to have a p_T between 30 GeV and 80 GeV. About 60% of events from LFV decays of the Higgs boson to $e\tau_\mu$ have a transverse momentum below 30 GeV as it can be seen from Figure 8.2(b). Since the muon of the WW +jets process is from a prompt decay of the W boson instead from a tau decay, the visible-mass has to fulfil a tighter requirement of $M_{\text{vis}} > 20$ GeV. Both leptons are required to have an angular distance in ϕ of $1.5 < \Delta\phi(e, \mu) < 2.8$ to suppress other background processes and to enrich the control region with WW +jets events. Events which have no jet are selected. A requirement on the angular distance between the muon

flying-direction and the $\vec{E}_T$ direction of $\Delta\phi(\mu, \vec{E}_T) > 1$ is applied, since $\sim 85\%$ of the LFV Higgs-boson decays to $e\tau_\mu$ have a $\Delta\phi$ smaller than 1, as shown in Figure 8.2(c). Finally, requirements on the transverse mass are applied: $m_T(e, \vec{E}_T) > 20$ GeV and $m_T(\mu, \vec{E}_T) > 30$ GeV. The tighter requirement on $m_T(\mu, \vec{E}_T)$ suppresses contributions from the signal process.

Events for the $WW \rightarrow 2l2\nu$ process were generated using POWHEG for the matrix element generation and PYTHIA 8 for the parton shower as presented in Table B.1 from Appendix B. The production cross-section for the inclusive WW+jets process was calculated including NNLO QCD effects [130] and found to be $118.7^{+2.5\%}_{-2.2\%}$ pb.

Furthermore, the W^+W^- cross section was measured in proton-proton collision at $\sqrt{s} = 13$ TeV with the CMS experiment:

$$\sigma_{W^+W^-} = 115.3 \pm 5.8 \text{ (stat)} \pm 5.7 \text{ (exp)} \pm 6.4 \text{ (theo)} \pm 3.6 \text{ (lumi)} \text{ pb [131].}$$

The total relative uncertainty on the measurement is 9% without considering the uncertainty on the luminosity. Predicted and measured production cross-sections are compatible with each other within the uncertainties. In the following the predicted cross section with a normalisation uncertainty of 10% is used for W^+W^- production.

Figure 8.18 shows the event yield in the W^+W^- process control region as function of $m_T(e, \vec{E}_T)$, and $m_T(\mu, \vec{E}_T)$. Predicted event yields are compatible within the uncertainties. Figure D.3 of Appendix D illustrates the event yield of the W^+W^- process for each jet category, applying all control region criteria except the requirement on the number of jets. An extra normalisation uncertainty of 5% is added to account for differences between the jet categories for the production cross-section..

WZ process

A control region for the WZ process is defined by requiring one isolated electron and two isolated muons, which have opposite charge. Events with further isolated leptons are vetoed. Requirements on the transverse momentum of the electron and the muon with largest p_T are $p_T^e > 30$ GeV and $p_T^\mu > 20$ GeV, respectively. Events which are compatible to have a Z boson decaying to a muon pair are selected, requiring an invariant dimuon-mass of $80 \text{ GeV} < M(\mu\mu) < 100 \text{ GeV}$. A veto on events with jets identified as b-quark jets is applied to exclude normalisation effects from the misidentification of jets as b-quark jets. Only events without jets are considered in order to suppress other background processes.

Events were generated with PYTHIA 8. Cross sections for the production of the WZ process were centrally computed using *Monte Carlo for FeMtobarn processes* (MCFM) [132], which is a program to calculate cross sections for various processes at hadron-hadron colliders with cross sections at the femtobarn-level. The matrix element for the inclusive WZ process is estimated at NLO accuracy. The computed cross section is 47.13 pb. The production cross-section [133] measured in proton-proton collisions at $\sqrt{s} = 13$ TeV is

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

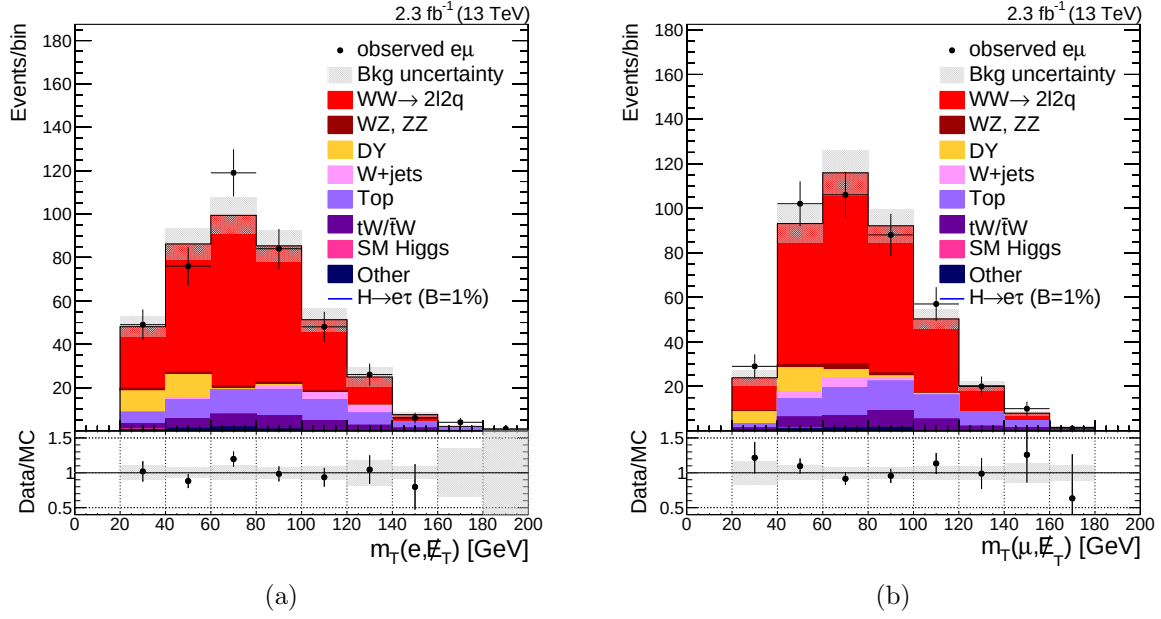


Figure 8.18.: Validation of the W^+W^- -process normalisation in the WW -process control-region. The event yield is shown as function of (a) $m_T(e, \cancel{E}_T)$, and (b) $m_T(\mu, \cancel{E}_T)$. A calculated inclusive cross section was used for W^+W^- production [130].

$\sigma(pp \rightarrow WZ) = 40.9 \pm 3.4$ (stat) $^{+3.1}_{-3.3}$ (syst) ± 0.4 (theo) ± 1.3 (lumi) pb. This measurement has a total uncertainty of $\sim 12\%$ without including the luminosity uncertainty. Predicted and measured values for the production cross-section differ by 15%.

Figure 8.19 shows the event yield in the WZ control-region as function of invariant dimuon-mass and jet multiplicity. Estimated and observed event yields are compatible within the uncertainties. There are deviations larger than the corresponding uncertainties for events with at least one jet, which increase with jet multiplicity. These deviations are negligible for the analysis because the top background dominates the 2-jet category. Diboson processes contribute to the 0-jet category and to a smaller amount to the 1-jet category [5].

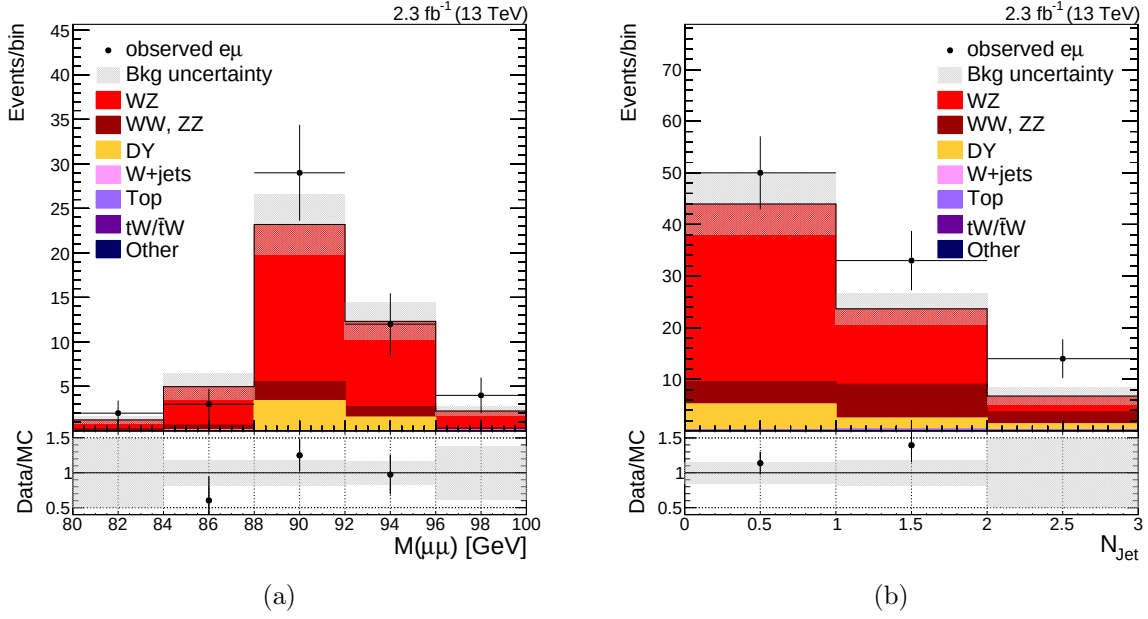


Figure 8.19.: Validation of the WZ-process normalisation in the WZ control-region. The event yield is shown as function of (a) invariant dimuon-mass, and (b) jet multiplicity after relaxing the requirement on the jet multiplicity to $N_{\text{Jet}} < 3$.

ZZ

A control region for the ZZ process can be defined by requiring two isolated electrons and two isolated muons. Events with further isolated leptons are vetoed. Only events with less than three jets are considered and a veto on jets identified as b-quark jets is applied. Events were generated with PYTHIA 8 and the cross section for the production of the ZZ process was centrally computed with NLO accuracy using MCFM [132]: $\sigma(\text{pp} \rightarrow \text{ZZ}) = 16.523 \text{ pb}$. The production cross-section of two Z bosons [134] was measured in proton-proton collisions at $\sqrt{s} = 13 \text{ TeV}$ employing leptonic decays of both Z bosons in the mass region $60 \text{ GeV} < m_Z < 120 \text{ GeV}$:

$$\sigma(\text{pp} \rightarrow \text{ZZ}) = 16.7_{-2.6}^{+2.9} (\text{stat})_{-0.5}^{+0.7} (\text{syst}) \pm 0.3 (\text{theo}) \pm 0.8 (\text{lumi}) \text{ pb}.$$

This measurement has a relative total-uncertainty of $\sim 18\%$ without including the luminosity measurement. Predicted and measured values for the ZZ production cross-section are compatible within the uncertainties and they only differ by 1%.

Figure 8.20 shows the event yield in the ZZ control region as function of invariant dielectron mass (a), invariant dimuon mass (b), and jet multiplicity (c). More collision data is needed to compare the predicted event and measured yield with each other as function of the invariant mass of the Z bosons. Event yields for events with one or two jets are compatible with the measurements within a normalisation uncertainty of 20%, while the normalisation uncertainty for events with no jets is $\sim 40\%$. The uncertainty for ZZ events with no jets is much higher than for the other diboson processes

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

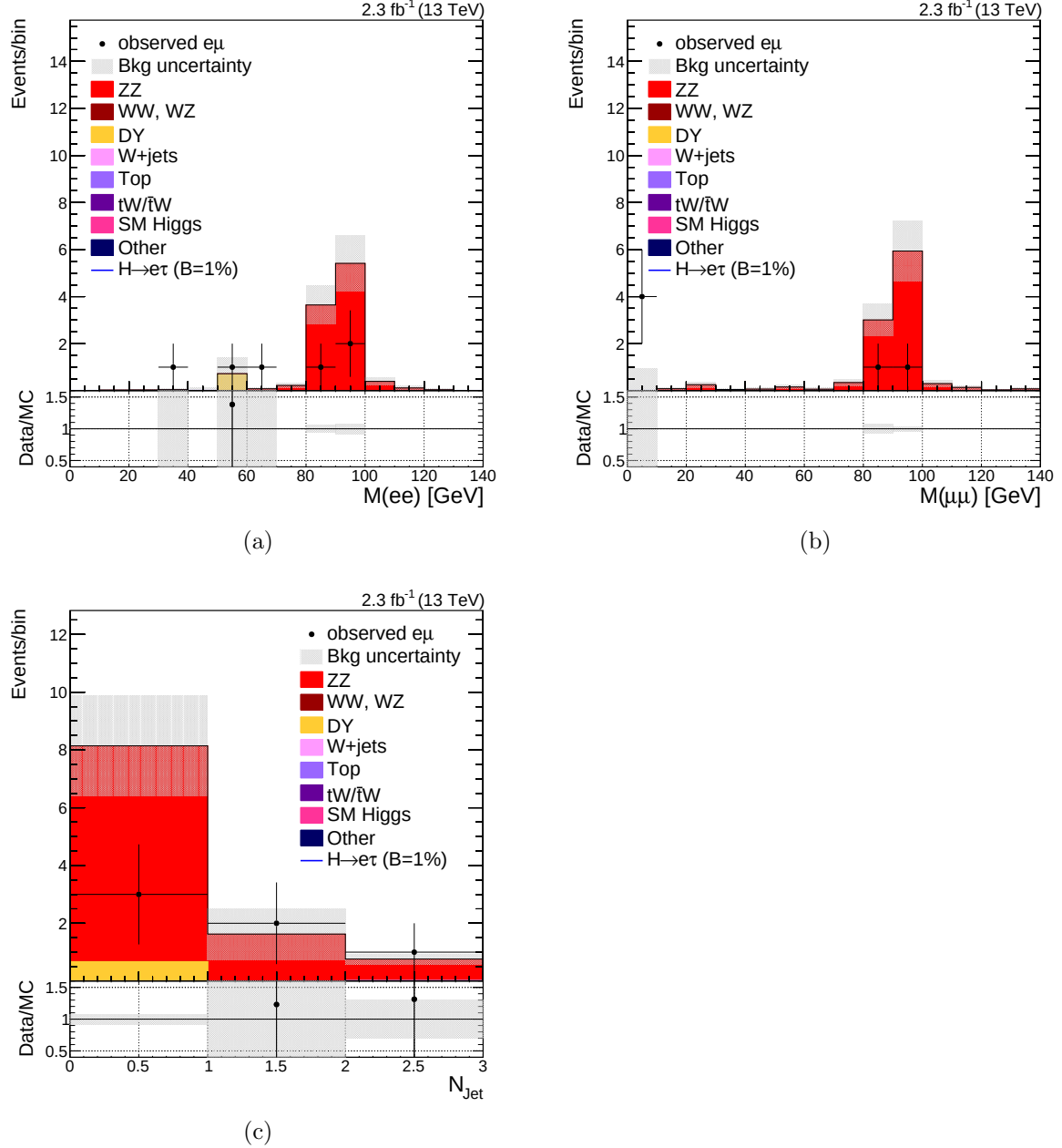


Figure 8.20.: Validation of the ZZ-process normalisation in the ZZ control-region. The event yield is shown as function of (a) invariant dielectron mass, (b) invariant dimuon mass, and (c) jet multiplicity. Events were generated with PYTHIA 8 and a computed inclusive ZZ cross-section was used.

Composition of the SM diboson background contribution

The SM diboson background is consisting of the $WW \rightarrow 2l2q$ process, the WZ process and the ZZ process. Figure 8.21 shows the composition of the diboson background as function of the jet multiplicity after preselection, and as function of the collinear mass for the 0-jet category. Compositions of the SM diboson prediction as function of the collinear mass are shown in Figure D.4 from Appendix D. The WW process dominates the other two diboson processes in all jet categories. This is also the case for each bin of the collinear mass. From this it follows that the normalisation uncertainty of the diboson process is in the order of the normalisation uncertainty of the WW production process and contributions from other diboson processes to the normalisation uncertainty can be neglected. In the following, a normalisation uncertainty of 10% is used for the diboson background with an additional 5% uncertainty for each jet category. These are the uncertainties estimated for the WW production process.

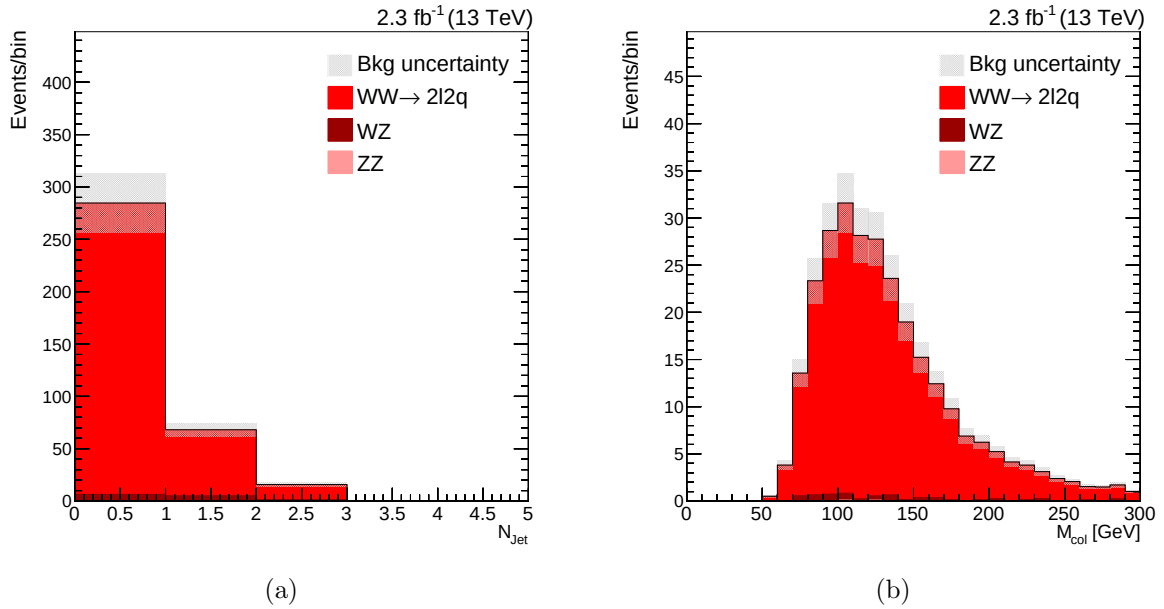


Figure 8.21.: Composition of the diboson background after preselection as function of (a) jet multiplicity, and (b) collinear mass in the 0-jet category. An uncertainty of 10% was used for each diboson process.

8.6.4. W+jets

The modelling of the misidentified lepton background was discussed in Section 8.5.2. A control region for the W+jets process with one misidentified lepton was introduced, where the two leptons are isolated and have the same charge. This control region is called region II. In the following, the normalisation of the misidentified lepton background is discussed. Figure 8.22 shows the event yield after preselection in this control region. It is shown as function of the collinear mass, and the jet multiplicity using a normalisation uncertainty of 40% for the misidentified-lepton background. Predicted and observed event yields as function of the collinear mass are compatible within background uncertainties. Simulation of the misidentified-lepton background underestimates the event yield, especially for events with a collinear mass of $50 \text{ GeV} < M_{\text{col}} < 100 \text{ GeV}$. Additionally, there is a deviation from the estimated yield for events with jets, which increases as function of the jet multiplicity. In Figure D.5 from Appendix D, background uncertainties were calculated using different normalisation uncertainties for the misidentified-lepton background. In the following, a normalisation uncertainty of 40% is used with an additional uncorrelated uncertainties of 10% for the 0-jet category, 40% for the 1-jet category, and 60% for the 2-jet category.

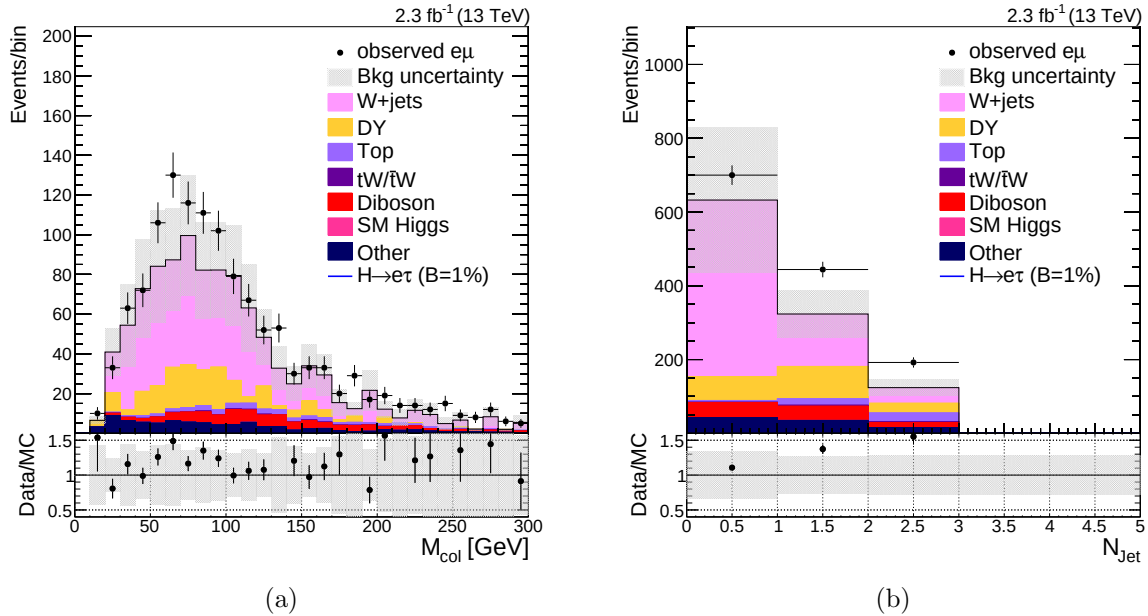


Figure 8.22.: Validation of the W+jets process normalisation in the same-charge region, region II. The event yield is shown after preselection as function of (a) collinear mass, and (b) jet multiplicity. A normalisation uncertainty of 40% was used for the misidentified-lepton background.

8.7. Profile likelihood method

A profile likelihood method is used to estimate the results on the branching fraction of the LFV Higgs-boson decay to $e\tau$. In the previous sections, the signal and background processes and their modelling were described. Production cross-sections of the SM Higgs boson are used for the signal model, while the branching fraction of the Higgs boson to $e\tau$ remains a free parameter. Signal and background model can be compared to collision data to discover the signal process, or to set limits on the branching fraction.

A null hypothesis H_0 can be defined, describing only known processes, which is tested against the alternative H_1 . For discovering a new signal process, the null hypothesis is the background-only hypothesis and the H_1 includes both, background and signal processes, while for setting limits, the H_0 is the model with signal plus background. Results on the branching fraction can be derived by using a likelihood ratio. In addition to parameters of interest, in this case the branching fraction, signal and background model contain in general *nuisance parameters* whose values must be derived from collision data [135].

The profile likelihood method is used to estimate the results on the branching fraction assuming the asymptotic approximation [136]. Distributions of the collinear mass for signal and various background processes are fitted to collision data. Systematic uncertainties are represented as nuisance parameters in the fit, and they can affect the signal and background normalisation, or also the shape of the collinear mass distribution. The CL_s method [137, 138] is used to set upper bounds on the branching fraction.

8.8. Systematic uncertainties

Systematic uncertainties are represented as nuisance parameter in the profile-likelihood fit. Some of these nuisance parameters affect the normalisation of all processes, while others only affect the normalisation of one process. There are also uncertainties which affect the shape of the collinear mass distribution. First, normalisation uncertainties are described which affect all processes. Then, normalisation uncertainties of background processes, followed by the normalisation uncertainties of the signal process, are discussed. Finally, uncertainties which have an effect on the collinear mass distribution are presented.

8.8.1. Normalisation uncertainties

Process independent normalisation uncertainties are listed in Table 8.7. The luminosity for the 2015 LHC proton-proton run at $\sqrt{s} = 13$ TeV was measured performing Van der Meer Scans [139], while the electron trigger, identification and isolation uncertainties can be estimated using the tag-and-probe technique [125]. The same method was used to estimate the muon identification and isolation efficiencies [125]. A systematic uncertainty

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

Table 8.7.: Systematic uncertainties in the expected event yield, which are treated as correlated between the jet categories. The uncertainty on the veto of jets identified as b-quark jets is applied on processes which do not include top-quarks.

Systematic uncertainty	
Luminosity	2.7%
Electron trigger/ID/isolation	3%
Muon ID/isolation	3%
Veto on jets identified as b-quark jets	3%

for jets misidentified as b-quark jets is considered for processes that do not include top-quarks [125], as a uncertainty on the identification efficiency for b-quark jets is included in the normalisation uncertainty of the top background.

8.8.2. Normalisation uncertainties of the background processes

In Section 8.6, control regions were defined for all background processes, except the $W\gamma$ process. It was discussed how systematic uncertainties can be derived using these control regions. Systematic uncertainties of various SM background processes are summarised in Table 8.8. The uncertainty of the $Z \rightarrow ll$ yield (10%) is derived from the measured [127] and from the predicted cross section. Uncertainties on the normalisation of the misidentified-lepton background are estimated by comparing the estimated event yield to collision data in the same-charge control-region. This uncertainty is in the order of 40%. As the measured cross section of the $t\bar{t}$ process differs from the predicted one by 10%, a larger systematic uncertainty is used in the analysis. This systematic uncertainty (20%) also accounts for the uncertainty on the identification efficiency of b-quark jets. There are no published measurements of the production cross section of the tW process for $\sqrt{s} = 13$ TeV, but it has been measured with $\sqrt{s} = 8$ TeV collision data (20% [140]). The diboson background is dominated by the W^+W^- process. Thus, its normalisation uncertainty (10%) is used for the diboson background. There is no published measurement of the cross section at $\sqrt{s} = 13$ TeV for the $W\gamma$ process, therefore the total uncertainty of the measurement of this cross section at $\sqrt{s} = 7$ TeV (10% [141]) is used as systematic uncertainty. This uncertainty also matches the systematic uncertainty on this background process of the LFV $\mu\tau$ search at $\sqrt{s} = 13$ TeV [125]. In addition to the normalisation uncertainty, a further systematic uncertainty is added for each background process, which is treated as uncorrelated between the jet categories. This uncertainty accounts for differences in the production cross-section between the jet categories. It can be estimated using the control regions for each background process. The value of 5% was taken from [125] and validated in the control regions. However, the misidentified lepton background has a higher uncertainty of 10% for the 0-jet category, 40% for the 1-jet category, and 60% for the 2-jet

Table 8.8.: Systematic uncertainties in the expected event yield for the jet categories. The first value gives the uncertainty, which is treated as correlated between the categories, and the second value is the uncertainty, which is treated as uncorrelated between jet categories.

Systematic uncertainty	0-jet	1-jet	2-jet
$Z \rightarrow ll$ background	$(10 \oplus 5)\%$	$(10 \oplus 5)\%$	$(10 \oplus 5)\%$
Misidentified-lepton background	$(40 \oplus 10)\%$	$(40 \oplus 40)\%$	$(40 \oplus 60)\%$
$t\bar{t}$ background	$(20 \oplus 5)\%$	$(20 \oplus 5)\%$	$(20 \oplus 5)\%$
tW background	$(20 \oplus 5)\%$	$(20 \oplus 5)\%$	$(20 \oplus 5)\%$
diboson background	$(10 \oplus 5)\%$	$(10 \oplus 5)\%$	$(10 \oplus 5)\%$
$W\gamma$ background	$(10 \oplus 5)\%$	$(10 \oplus 5)\%$	$(10 \oplus 5)\%$

category. The uncertainty of the 0-jet category matches the corresponding uncertainty from [125], while for the other jet categories higher uncertainties are found.

8.8.3. Normalisation uncertainties of the Higgs boson processes

There are theoretical uncertainties which affect the normalisation of the Higgs boson processes. Theoretical uncertainties on the production cross-section affect both, SM Higgs boson production and LFV Higgs-boson production. In the cross-section calculation of the Higgs boson, QCD processes were calculated with NNLO+NNLL accuracy for gluon fusion, and with NNLO accuracy for vector boson fusion. Electroweak processes were calculated with NLO accuracy [142]. The QCD scale was set to the mass of the Higgs boson $\mu = \mu_F = \mu_R = M_H$, with the factorisation scale μ_F , and the renormalisation scale μ_R . The QCD scale uncertainty was estimated by varying μ in the range $\frac{1}{2} < \frac{\mu}{M_H} < 2$ constraining $\frac{\mu_F}{\mu_R}$ to $\frac{1}{2} < \frac{\mu_F}{\mu_R} < 2$. Different parton distribution functions (PDFs) were combined to calculate the cross sections. Their uncertainties are used to estimate the PDF uncertainties. These uncertainties also include the uncertainty on the strong coupling-constant α_s . In addition, the systematic uncertainty on the branching fraction of the SM Higgs-boson decay to two taus is 1.17%. Table 8.9 lists the resulting uncertainties on the Higgs boson processes.

Table 8.9.: Theory uncertainties of the Higgs boson processes

Theory	QCD scale [%]	PDF [%]	combined
Gluon fusion	+7.6/ − 8.1	3.1	8.7%
Vector boson fusion	+0.4/ − 0.3	2.1	2.1%
$B(H \rightarrow \tau\tau)$ [%]	Uncertainty [%]		
6.272	1.17		

8.8.4. Shape uncertainties of the collinear mass

Some nuisance parameter also affect the shape of the collinear mass distribution. Corrections are applied for the jet energy scale (JES) as function of jet p_T and η . They are propagated to the missing transverse energy. Uncertainties of the JES affect the normalisation and the shape of the collinear mass distribution for signal and background processes. Shape and normalisation uncertainties related to the JES are computed by altering the JES by one standard deviation $\pm\sigma$. The resulting collinear mass distributions are used in the fit. There is also an uncertainty to account for the unclustered energy scale. Unclustered energy comes from jets with $p_T < 10$ GeV and from PF candidates which were not clustered to any jet. Shape uncertainties are calculated by altering the unclustered energy scale by $\pm\sigma$ and propagating changes to $\cancel{E}_T$. These effects cause a shift of the collinear mass distribution and the resulting distributions are used in the fit.

8.9. Discriminating variables

In this section, variables which have discriminating power to separate signal and background processes are introduced. Some of them are already used in the $\sqrt{s} = 8$ TeV analysis, while others have not been used previously in searches for LFV decays of the Higgs boson at the CMS experiment. The collinear mass distributions are shown blinded in the signal region in order not to bias the result. All distributions in this section are shown before the fit.

First, the collinear mass distributions are shown after preselection. Then, the discriminating variables, which were used in the $\sqrt{s} = 8$ TeV analysis are discussed. Finally, the new discriminating variables are introduced.

8.9.1. Collinear mass distributions after preselection

The collinear mass, described in Section 8.2, is used to define the signal region. LFV decays of the Higgs boson have a collinear mass with a mean value of 125 GeV and a width in the order of 15 GeV, therefore the signal region is defined in the mass window $100 \text{ GeV} < M_{\text{col}} < 150 \text{ GeV}$. In the fit, the other bins of the collinear mass distributions are used to fit the nuisance parameters, i.e. the normalisation of the background contributions.

Figure 8.23 shows the collinear mass for each jet category. The signal has its peak between $100 \text{ GeV} < M_{\text{col}} < 150 \text{ GeV}$, and the agreement between estimated and observed background expectation is good.

The signal region of the 0-jet category is dominated by the Drell-Yan and the diboson background contributions. Misidentified-lepton background and the top background also contribute. There is a Z-boson peak coming from the Drell-Yan background contribution below 100 GeV. This peak can be used in the fit to estimate the Drell-Yan normalisation. The side band of $150 \text{ GeV} < M_{\text{col}} < 300 \text{ GeV}$ can be used to constrain the diboson and the top background contributions. For the 1-jet and 2-jet categories, the top background is the dominant background process with additional contributions from the Drell-Yan background, and the diboson processes. The fraction of the misidentified-lepton to the total expected background contribution decreases with the jet multiplicity and is negligible for the categories which include at least one jet compared to the other background contributions.

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

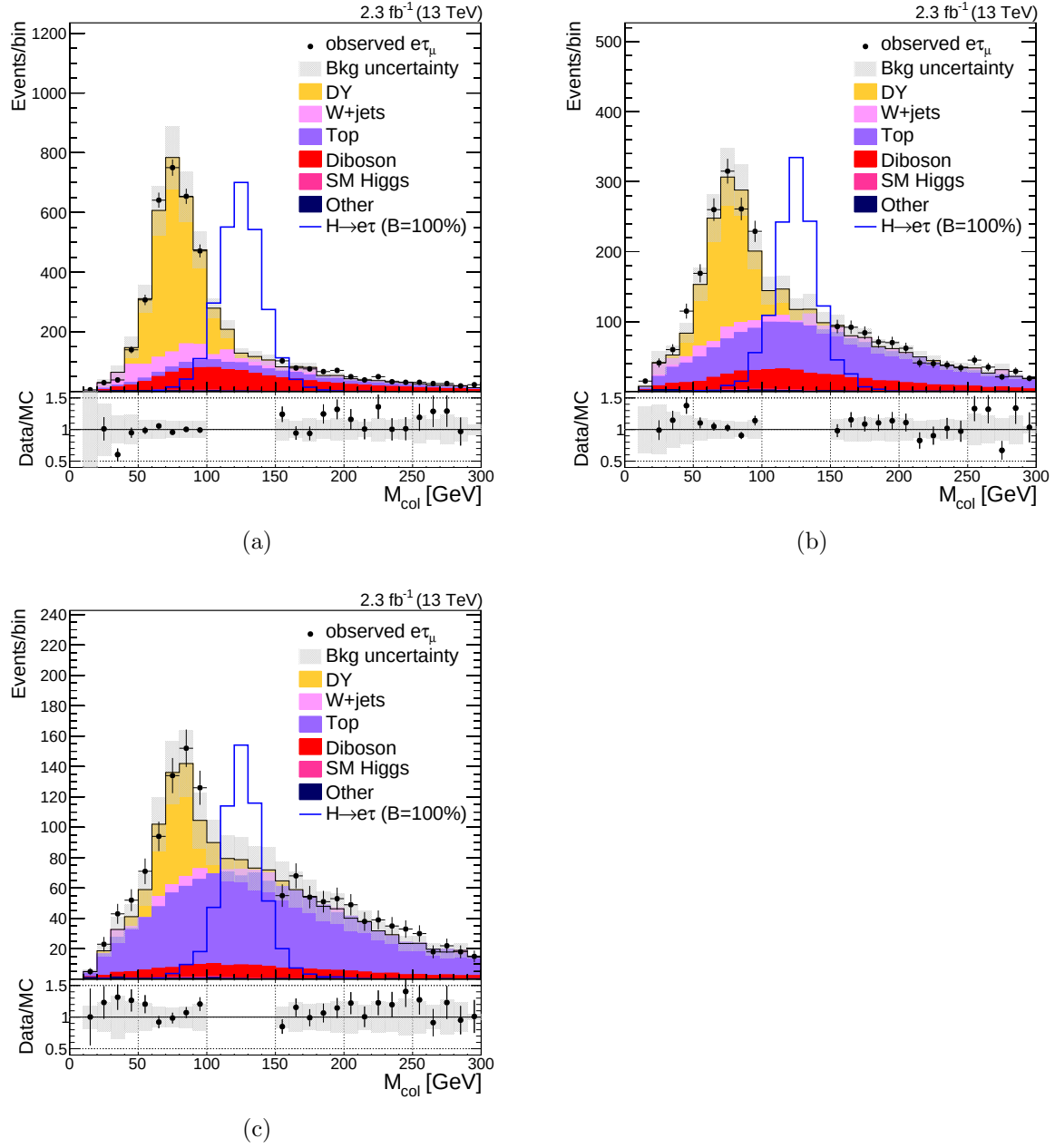


Figure 8.23.: Collinear mass after preselection for the 0-jet category (a), 1-jet category (b), and 2-jet category (c) for proton-proton collisions at $\sqrt{s} = 13$ TeV. Distributions are shown blinded in the signal region.

8.9.2. Discriminating variables used in the $\sqrt{s} = 8$ TeV analysis

There are several discriminating variables which were used in the search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$ at $\sqrt{s} = 8$ TeV [5], and in the first direct search for LFV decays of the Higgs boson [17]. These variables are kinematic variables, such as transverse momenta, angular variables, and mass variables. In the following, all distributions are shown after the preselection, which is discussed in Section 8.4.4. First, transverse momenta of leptons are described. Then, angular variables are discussed followed by the transverse mass between the muon flying direction and $\vec{E}_T$. Finally, variables are introduced which are sensitive to VBF production.

The transverse momentum of the electron (p_T^e) is used to preselect events which pass the single electron trigger. As described in Section 8.1, the transverse momentum of the electron can be used to distinguish between SM Higgs-boson decays to two taus and the LFV Higgs-boson decay to $e\tau$, due to the harder momentum spectrum of the prompt electron with respect to the electron from the tau decay. Figure 8.24 shows the transverse momentum of the electron for each jet category. Signal contributions peak between $50 \text{ GeV} < p_T^e < 65 \text{ GeV}$ for the 0-jet category. A tighter criterion can reduce the amount of contributions from the dominant Drell-Yan background, the misidentified-lepton background, and the diboson processes. The distributions of the transverse momentum for the other categories are broader than for the 0-jet category. Some background contributions could be reduced by a tighter criterion on the electron transverse momentum for the 1-jet category, but the distribution of the 2-jet category is almost flat.

The transverse momentum of the muon (p_T^μ) has a softer spectrum than the transverse momentum of the electron, as the muon is coming from the tau decay. For all jet categories, the spectrum is falling as it can be seen in Figure 8.25. A tighter criterion on the transverse momentum of the muon would reduce the amount of misidentified leptons to the background contributions significantly. It would also reduce the background contribution from Drell-Yan events.

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

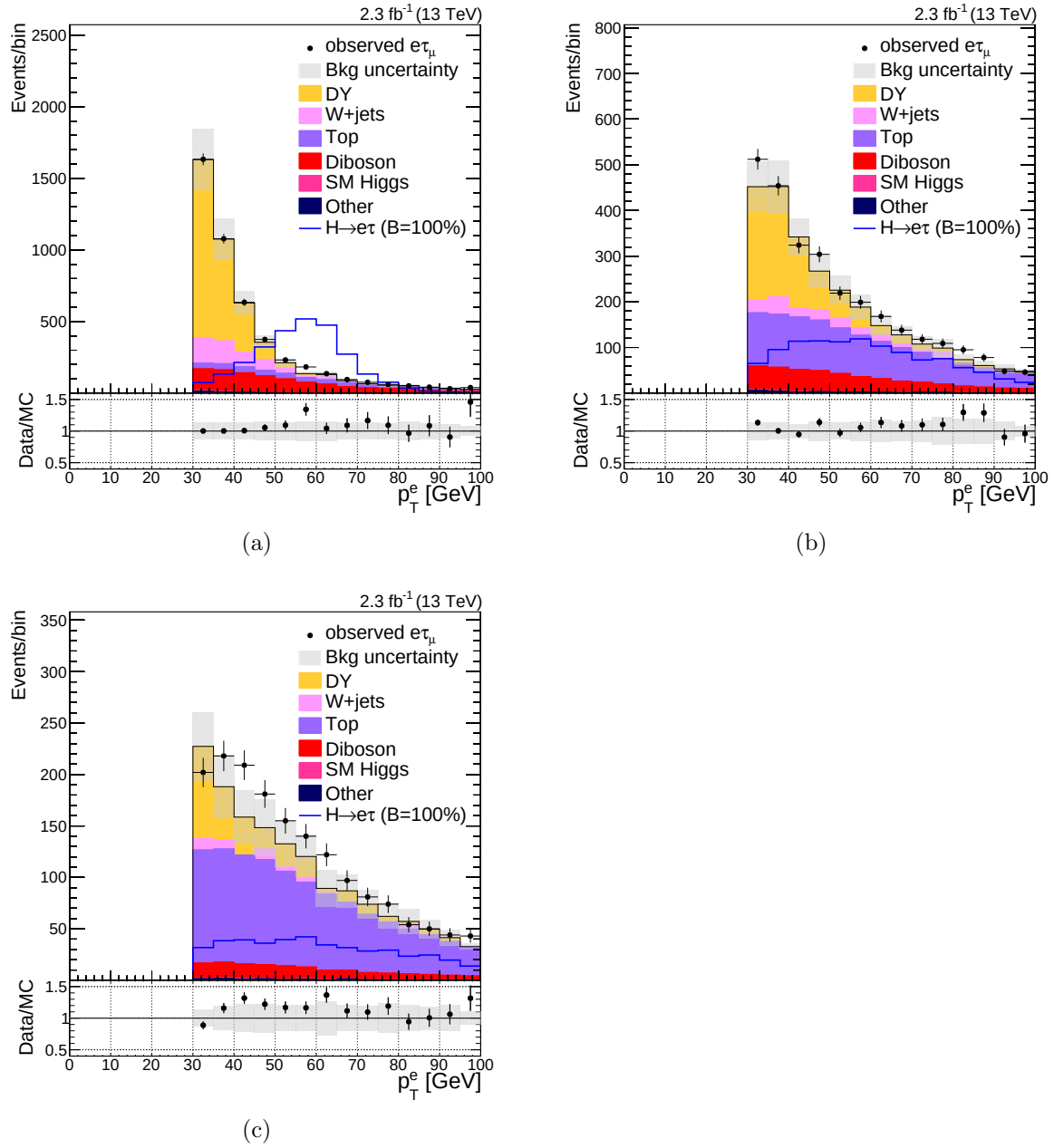


Figure 8.24.: Transverse momentum of the electron after preselection for the 0-jet category (a), 1-jet category (b), and 2-jet category (c) for proton-proton collisions at $\sqrt{s} = 13$ TeV.

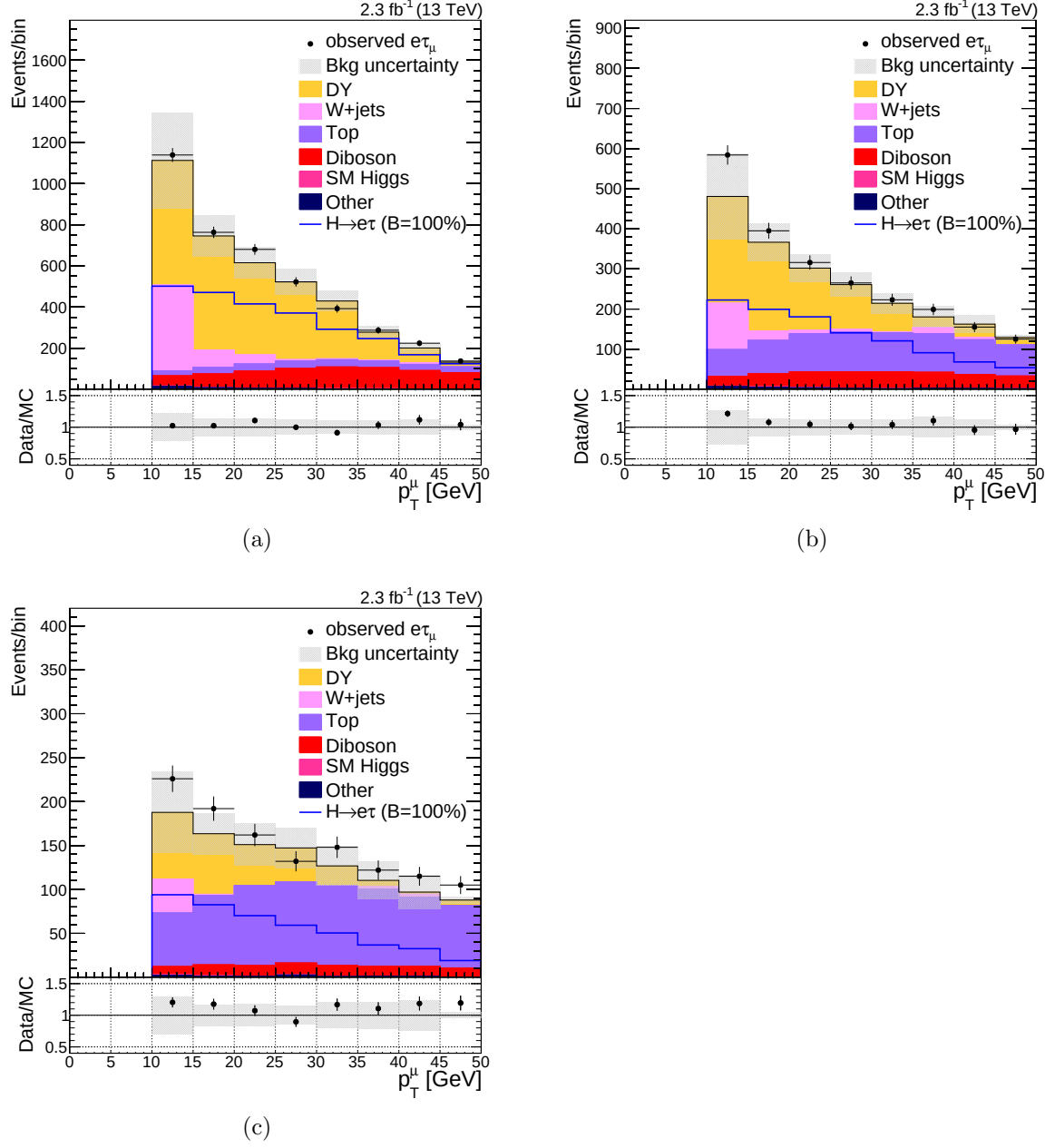


Figure 8.25.: Transverse momentum of the muon after preselection for the 0-jet category (a), 1-jet category (b), and 2-jet category (c) for proton-proton collisions at $\sqrt{s} = 13$ TeV.

Angular discriminating variables which were used in the $\sqrt{s} = 8$ TeV analysis are the azimuthal angle between the muon flying-direction and $\vec{E}_T$, denoted as $\Delta\phi(\mu, \vec{E}_T)$, and the azimuth angle between the electron and muon flight-directions, $\Delta\phi(e, \mu)$. Figure 8.26 shows the event yield as function of both discriminating variables for the 1-jet category. The decay products of the tau are collinear to the tau flight-direction, due to the large mass difference between the Higgs boson and the tau. Thus, $\Delta\phi(\mu, \vec{E}_T)$ peaks at 0 for the signal process, while the distribution of the total background expectation is flat as function of $\Delta\phi(\mu, \vec{E}_T)$. Both leptons tend to have flight directions in the opposite directions because the Higgs boson decays to an electron and a tau, which further decays to a muon with the same flying direction as the tau. The distribution of the background contributions is similar to the one of the signal process, but a criterion on $\Delta\phi(e, \mu)$ can be used to get reduced background contributions in the tail of $\Delta\phi(e, \mu)$ without losing much signal events.

The transverse mass between the muon flying-direction and $\vec{E}_T$, $m_T(\mu, \vec{E}_T)$, is defined as:

$$m_T(\mu, \vec{E}_T) = \sqrt{2 p_T^\mu E_T [1 - \cos(\Delta\phi)]}, \quad (8.7)$$

with $\Delta\phi$ being the angular distance in ϕ between the muon flying-direction and the missing transverse energy vector. Figure 8.26 shows the event yield as function of $m_T(\mu, \vec{E}_T)$ for the 1-jet category. Since $\Delta\phi(\mu, \vec{E}_T)$ peaks around 0, the transverse mass also peaks at 0. Contributions from the background processes tend to have larger values than signal contributions. Especially the top background has a peak in a regions of the transverse mass distribution where the signal contributions are negligible, and it can be significantly reduced by applying an upper requirement on $m_T(\mu, \vec{E}_T)$.

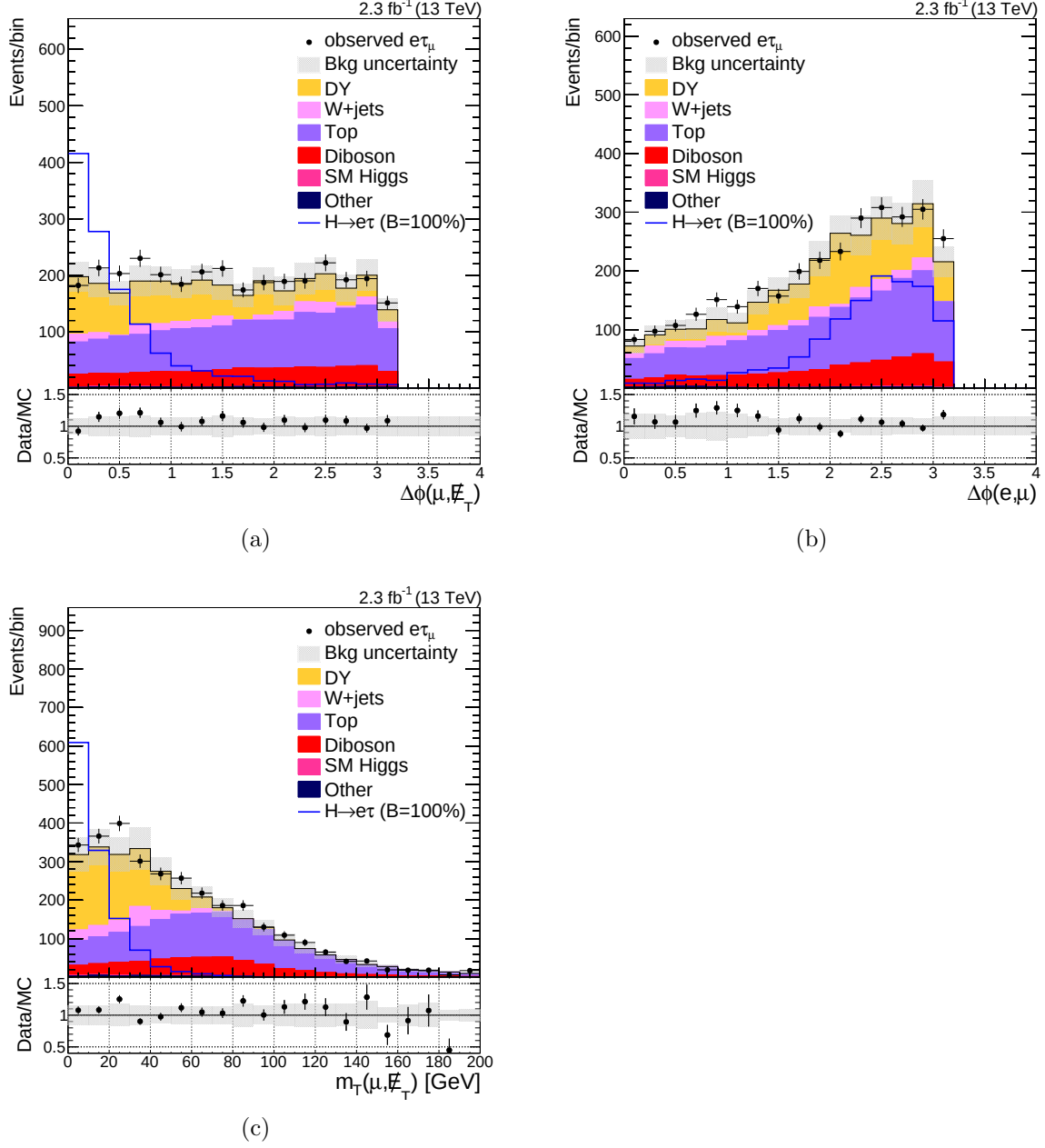


Figure 8.26.: Distributions of discriminating variables used in the $\sqrt{s} = 8$ TeV analysis. The event yield is shown as function of $\Delta\phi(\mu, \cancel{E}_T)$ (a), $\Delta\phi(e, \mu)$ (b), and $m_T(\mu, \cancel{E}_T)$ (c) for the 1-jet category.

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

VBF variables are discriminating variables that select events from vector boson fusion. Such variables are the distance in η between two jets, $|\Delta\eta|$, and the invariant dijet mass M_{jj} . Figure 8.27 shows the distribution of both variables. Background contributions of the 2-jet category peak at smaller values of $|\Delta\eta|$, while the signal contributions have a flat-falling spectrum. The invariant dijet mass spectrum peaks around 100 GeV for signal and background processes, with a steeper falling background spectrum. Tight criteria on both variables would enhance the fraction of VBF events, but they would also reduce events from LFV Higgs-boson decays to $e\tau$, where the Higgs boson was produced via gluon fusion. About half of the events in the 2-jet category are produced via gluon fusion. The other half of the signal events is produced via VBF. Thus, a tight VBF selection would reduce the amount of signal, and therefore the sensitivity of the analysis. An extra category for VBF events could increase the sensitivity, as the 2-jet category could be split into two categories with dedicated selections on the Higgs-boson production they are sensitive to. This is not possible for the 2015 collision data because of the small statistic in the 2-jet category.

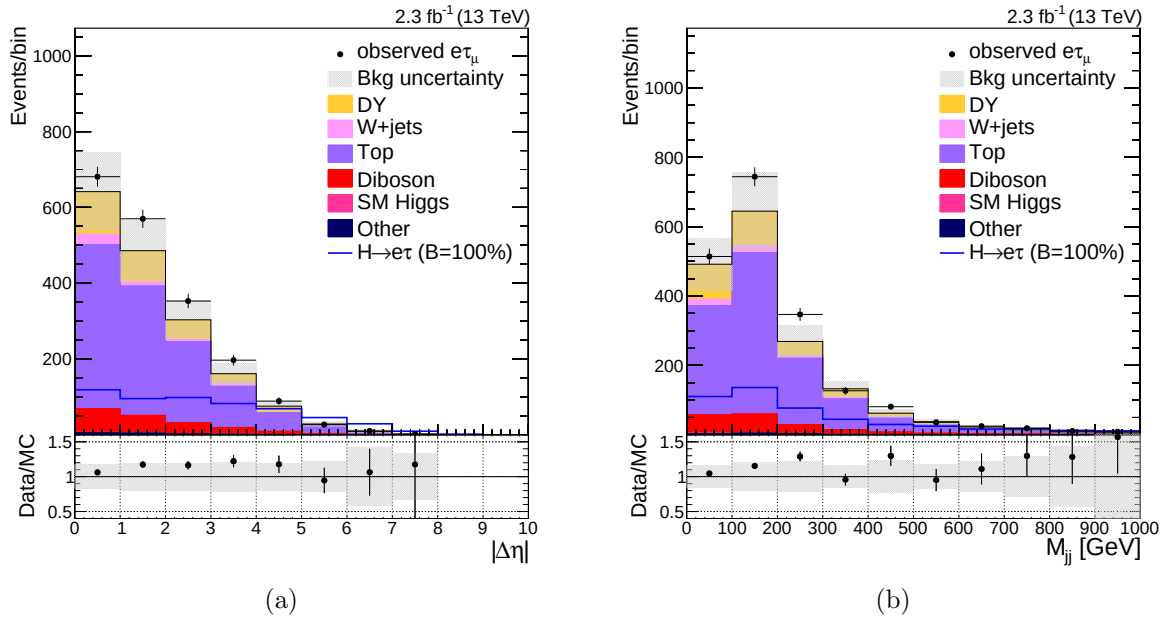


Figure 8.27.: Discriminating VBF variables of the $\sqrt{s} = 8$ TeV analysis. The event yield is shown as function of $|\Delta\eta|$ (a), and M_{jj} (b) for the 2-jet category.

8.9.3. New discriminating variables

New discriminating variables were investigated for the search at $\sqrt{s} = 13$ TeV. In the mass reconstruction of the Higgs boson, presented in Section 8.2, the neutrino momentum is estimated by the part of the transverse missing energy which is collinear to the direction of the muon flying-direction, $\vec{E}_{T,\text{proj}}$. If the $\vec{E}_T$ is pointing to the opposite direction with respect to the muon flying-direction, $\vec{E}_{T,\text{proj}}$ is negative. Neutrinos originating from the $H \rightarrow e\tau$ process are in most cases flying in the same direction as the muon, therefore $\vec{E}_{T,\text{proj}}$ has to be positive for this events. Figure 8.28 shows the event yield as function of the new discriminating variables. The signal distribution peaks at positive values for $\vec{E}_{T,\text{proj}}$ and has only a small tail to negative values, while the background peaks at zero and is similarly distributed for negative and positive $\vec{E}_{T,\text{proj}}$ values. However, the discriminating variable $\Delta\phi(\mu, \vec{E}_T)$ is strongly correlated with the sign and amplitude of $\vec{E}_{T,\text{proj}}$, while $m_T(\mu, \vec{E}_T)$ is correlated with the amplitude of $\vec{E}_{T,\text{proj}}$. Thus, only a small impact on the expected upper limit is expected, when using $\vec{E}_{T,\text{proj}}$ as additional discriminating variable. This variable can be used to define a cleaner signal region by only considering events where $\vec{E}_{T,\text{proj}}$ and the muon-flying direction point in the same direction.

Another discriminating variable, which is used in the collinear mass definition, is the visible energy fraction of the tau decay-products, x_τ^{vis} . The muons from the signal process carry only a part of the tau energy. It can be seen in Figure 8.2, that the signal distribution of x_τ^{vis} has its peak below 0.5, while the background peaks at higher values. This variable has additional discriminating power with respect to $\Delta\phi(\mu, \vec{E}_T)$ and $m_T(\mu, \vec{E}_T)$, as it combines information of the muon p_T relative to sum of p_T^μ and the absolute value of $\vec{E}_T$.

The discriminating variable A_ϕ combines the angular information between the electron flying-direction, the muon flying-direction, and $\vec{E}_T$. A_ϕ is defined as:

$$A_\phi = \frac{\Delta\phi(e, \vec{E}_T) - \Delta\phi(\mu, \vec{E}_T)}{\Delta\phi(e, \mu)}, \quad (8.8)$$

with $\Delta\phi(e, \vec{E}_T)$ being the azimuthal distance between the electron flying-direction and $\vec{E}_T$. The distribution of $\Delta\phi(e, \mu)$ peaks at π for the signal process, as the electron and muon have a higher probability to have opposite flying directions. Since the neutrinos are collinear to the muon, $\Delta\phi(\mu, \vec{E}_T)$ tends to be very small and $\Delta\phi(e, \vec{E}_T)$ is approximately as large as $\Delta\phi(e, \mu)$. Thus, the signal-process distribution of A_ϕ peaks at 1. The background distribution also peaks at $A_\phi = 1$, but it has a larger width to smaller values.

The last new introduced variable is $x_{\vec{E}_T}^\nu$, which is defined as fraction of $\vec{E}_{T,\text{proj}}$ to the full $\vec{E}_T$:

$$x_{\vec{E}_T}^\nu = \frac{\vec{E}_{T,\text{proj}}}{\vec{E}_T}. \quad (8.9)$$

If the missing transverse energy originates from neutrinos which are collinear to the

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

muon-flying direction, $\cancel{E}_{T,\text{proj}}$ and $\cancel{E}_T$ have similar values. For events with several neutrinos from different decays, $\cancel{E}_T$ is the absolute value of the vector sum of the neutrino transverse-momenta and $x_{\cancel{E}_T}^\nu$ tends to have smaller values. Events with no genuine $\cancel{E}_T$ have arbitrary values of $\cancel{E}_{T,\text{proj}}$, therefore the distribution of $x_{\cancel{E}_T}^\nu$ for such events should differ to the signal process one. It can be seen from Figure 8.28 that the signal distribution peaks at $x_{\cancel{E}_T}^\nu = 1$, while the background peaks at $x_{\cancel{E}_T}^\nu = 1$ and $x_{\cancel{E}_T}^\nu = -1$ with a flat distribution between both values. Thus, requiring $x_{\cancel{E}_T}^\nu$ to have positive values can reduce some background without losing much signal events.

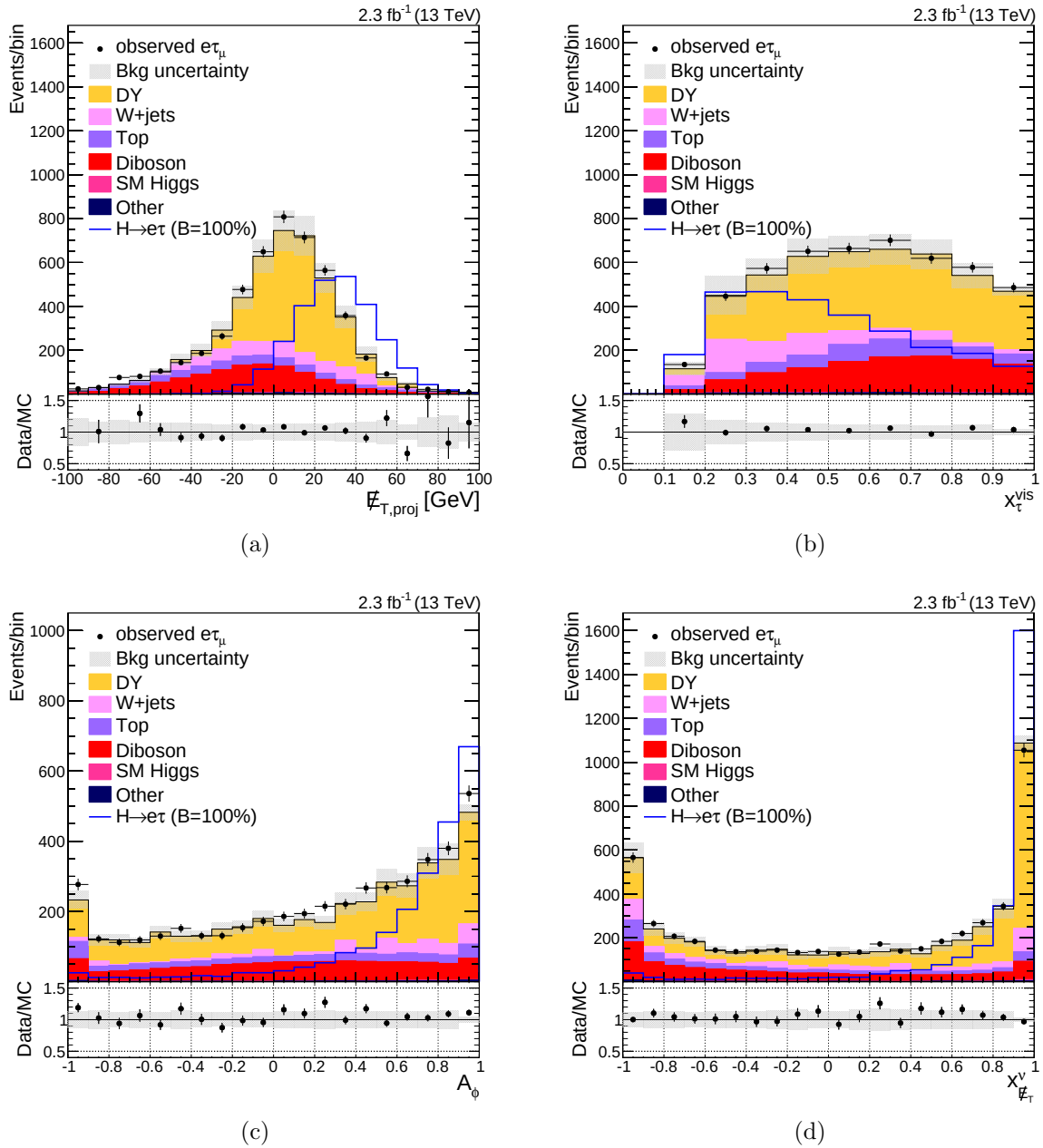


Figure 8.28.: New discriminating variables- The event yield is shown after preselection for the 0-jet category as function of $\cancel{E}_{T,\text{proj}}$ (a), $x_{\cancel{E}_T}^{\text{vis}}$ (b), A_ϕ (c), and $x_{\cancel{E}_T}^\nu$ (d).

8.10. Search for lepton-flavour violating Higgs-boson decays at $\sqrt{s} = 8$ TeV

The search for lepton-flavour violating decays of the Higgs boson to $e\tau$ at $\sqrt{s} = 8$ TeV was published in [5]. It is the first direct search for LFV decays of the Higgs boson to $e\tau$. In the following, the $e\tau_\mu$ channel is described. There are some differences in the used techniques for the $\sqrt{s} = 8$ TeV analysis with respect to the $\sqrt{s} = 13$ TeV one.

Event reconstruction: Electrons, muons, jets, and missing transverse energy are reconstructed as described in Section 6. Jets are reconstructed by clustering all PF candidates with the anti- k_t algorithm [112] using a distance parameter of $R = 0.5$, instead of $R = 0.4$ used at $\sqrt{s} = 13$ TeV. Jet energy scale corrections are applied and a pileup reweighting is done. PF $\cancel{E}_T$ is used in this analysis and Type-I corrections are applied. Muons are required to pass the tight muon selection, described in Section 6.3. A medium muon ID was not available, as it was introduced in LHC Run II. Electrons have to pass a tight criterion on a discriminator value of a BDT, which was trained for electrons at $\sqrt{s} = 8$ TeV.

Loose event selection: Events were collected in proton-proton collisions at $\sqrt{s} = 8$ TeV with a bunch spacing of 50 ns. They were recorded with the CMS detector in 2012 and correspond to an integrated luminosity of 19.7 fb^{-1} .

A muon-electron trigger, `HLT_Mu8_Ele17_CaloIdT_CaloIsoVL_TrkIdVL_TrkIsoVL`, was used to collect the collision data sample. Events were collected with at least one electron with $p_T^e > 17$ GeV and $|\eta^e| < 2.5$, and one muon with $p_T^\mu > 8$ GeV and $|\eta^\mu| < 2.4$. The trigger also applies loose identification and isolation requirements to the leptons.

An event cleaning, as presented in Section 8.4.4 for the $\sqrt{s} = 13$ TeV analysis, is also performed. An electron with $p_T^e > 40$ GeV and $|\eta^e| < 2.3$, excluding the transition region, and a muon with $p_T^\mu > 10$ GeV and $|\eta^\mu| < 2.1$ of opposite charge are required. Events which have additional leptons with $p_T > 7$ GeV are rejected. Jets have to fulfil the same identification requirements as for the $\sqrt{s} = 13$ TeV analysis. If an event has jets identified as b-quark jets, this event is not considered further. A selection on the visible mass of $M_{\text{vis}} > 12$ GeV is applied.

Trigger efficiency: In the $\sqrt{s} = 8$ TeV analysis, the MC samples are rescaled by the measured efficiencies instead of applying the trigger to simulated events. The efficiency of the trigger selection is measured from data. It is based on a single-muon and a single-electron collision data set which were collected using the reference triggers `HLT_IsoMu24_eta2p1` or `HLT_IsoMu24`, and `HLT_Ele27_WP80`, respectively. Efficiencies are calculated separately for muons and electrons from the ratio of the events firing the selection trigger divided by

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

the number of events firing the reference trigger. The product of the muon and electron trigger efficiencies gives the total trigger efficiency.

Background modelling: The same background processes contribute to the background as described in Section 8.3. All processes are taken from MC simulation, except the two dominant background processes, which are the $Z \rightarrow \tau\tau$ background, and the misidentified lepton background. Both background contributions are estimated using collision data as described in Section 8.5. The $Z \rightarrow \tau\tau$ background is estimated using an embedded sample. A shift of 2% in the mass peak relative to simulation is observed, since differences in the final-state radiations of photons between muons and taus are not taken into account. The isolation side-band method is used for estimating the misidentified-lepton background. Only misidentified muons are considered, as misidentified electrons and double fakes can be neglected. $Z \rightarrow \mu\mu + X$ events were used to measure the misidentification ratio in an independent data sample. Dimuon collision events were collected using muon triggers. Figure 8.29 shows the collinear mass after the loose selection in the same-charge control region, region II. The estimated and observed event yields agree within the uncertainties, including a systematic uncertainty of 40% for the misidentified lepton background. A $t\bar{t}$ enriched control region is defined by the loose selection for the 2-jet category. At least one jet identified as b-quark jet is required. The normalisation is checked using this control region. A difference between simulated and measured p_T spectrum of the top-quarks was seen [143], therefore the transverse momentum of the top-quark in simulation was reweighted accordingly. The collinear mass distributions after loose selection for each jet category are shown in Figure E.1 in Appendix E.

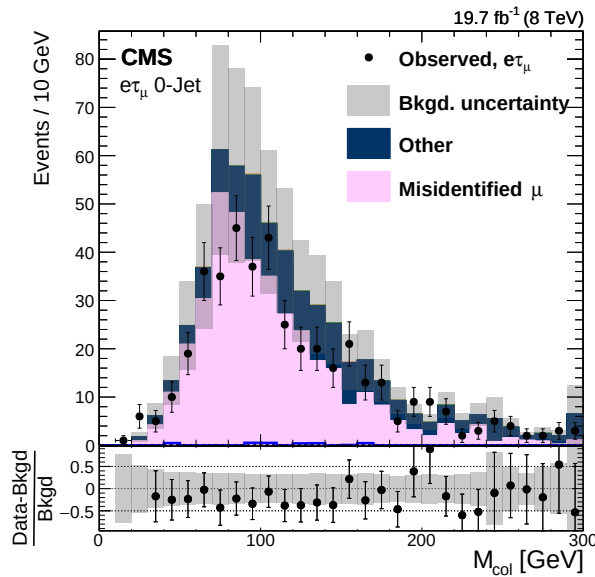


Figure 8.29.: Collinear mass distribution after applying a loose selection in region II, where both leptons have opposite charge. Events were collected in proton-proton collisions at $\sqrt{s} = 8$ TeV [5].

Event selection: The event selection was optimised to maximise the significance in the signal region:

$$sig = \frac{S}{\sqrt{S+B}}, \quad (8.10)$$

with the event yield S of the signal process and the event yield B of the background contributions. The sensitivity of the $\mu\tau_e$ channel of the search for $H \rightarrow \mu\tau$ analysis is in the same order of magnitude as for the $e\tau_\mu$ channel in the search for $H \rightarrow e\tau$ because only kinematic selections are varied. The upper limit on the branching fraction of $\mu\tau_e$ was in the order of 1% for the analysis at $\sqrt{s} = 8$ TeV [17], therefore $B(H \rightarrow e\tau) = 1\%$ was used in the optimisation of the event selection.

A new technique is used to optimise the event selection. First, a set of discriminating variables is defined. These variables are introduced in Section 8.9.2. A tighter requirement on the muon transverse momentum of $p_T^\mu > 15$ GeV is applied to reduce the misidentified lepton background. Then, the event selection is optimised for each jet category. The optimisation starts by scanning all discriminating variables in specified step sizes. Each variable is scanned individually by applying selection requirements on this variable and calculating the significance for the applied criterion. This is illustrated in Figure 8.30 for the discriminating variable $\Delta\phi(\mu, \cancel{E}_T)$. Two separate scans are done, one for each direction selecting events with larger/smaller values ($>/<$), i.e. the significance is calculated for the requirement $\Delta\phi(\mu, \cancel{E}_T) > 0.1$, then for $\Delta\phi(\mu, \cancel{E}_T) > 0.2$, and so on. The same is done for requiring smaller values. After both scans, the best selection requirement on a variable is determined which maximises the significance. In this case, the requirement $\Delta\phi(\mu, \cancel{E}_T) < 0.8$ gives the best significance. Finally, the selection criteria are ranked according to their significance and the selection which maximises the significance is applied. If there are two criteria with same significance, the criterion with larger signal event yield is applied. The selected variable is removed from the list of discriminating variables. If the requirement $\Delta\phi(\mu, \cancel{E}_T) < 0.8$ would have improved the event selection the most, this requirement would be applied and $\Delta\phi(\mu, \cancel{E}_T)$ would be removed from the list of discriminating variables. Then, the remaining discriminating variables are scanned again. This procedure stops, when the significance does not improve by applying a further selection, or when there are no further variables left.

The resulting event selection is listed in Table 8.10. Looser VBF selection criteria are applied with respect to the $\mu\tau$ analysis [17], as signal events from gluon fusion contribute to about 50% to the signal contribution to the 2-jet category. The sensitivity of this category can be improved by loosening the VBF criteria.

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

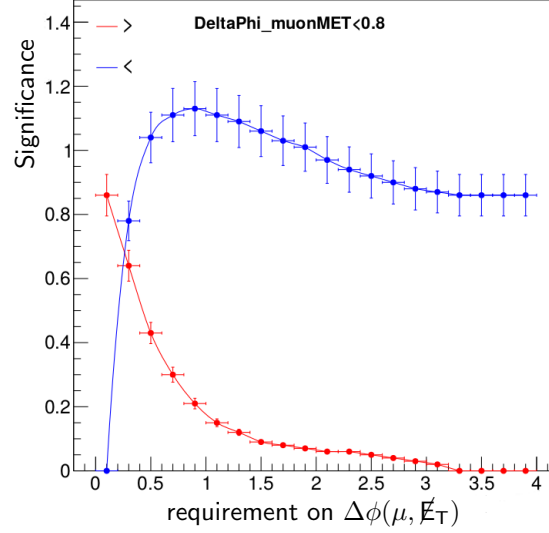


Figure 8.30.: Optimisation scan for significance for $\Delta\phi(\mu, \cancel{E}_T)$ in the 0-jet category for proton-proton collisions at $\sqrt{s} = 8$ TeV. A variable is scanned by applying selection criteria larger/smaller than the given value of the discriminating variable and calculating the significance for the applied criterion. The red curve shows the scan for requiring the events to have larger ($>$) values for this variable than the criterion, while the blue curve shows the scan, where the events are required to have smaller ($<$) values.

Table 8.10.: Event selection of search for LFV decays of the Higgs boson to $e\tau_\mu$ at $\sqrt{s} = 8$ TeV [5].

Variable	0-jet	1-jet	2-jet
p_T^e	> 50 GeV	> 40 GeV	> 40 GeV
p_T^μ	> 15 GeV	> 15 GeV	> 15 GeV
$m_T(\mu, \cancel{E}_T)$	-	< 30 GeV	< 40 GeV
$\Delta\phi(\mu, \cancel{E}_T)$	< 0.8	< 0.8	-
$\Delta\phi(e, \mu)$	-	> 0.5	-
$ \Delta\eta $	-	-	> 3
M_{jj}	-	-	> 200 GeV

Systematic uncertainties: A profile likelihood fit described in Section 8.7 is used. Systematic uncertainties contribute as nuisance parameters to the fit. These nuisance parameters can affect the background and signal normalisation, or they can also affect the shape of the collinear mass distributions. The peak of the collinear mass is shifted by 2% in the embedded sample. A correction for this is applied and the systematic uncertainty is evaluated by shifting the peak of the embedded sample by 2% in both directions and using the resulting collinear mass distributions in the fit. Shape and normalisation uncertainties related to the JES are computed by altering the JES by one standard deviation. The effect on the JES is also propagated to $\cancel{E}_T$. The resulting collinear mass distribution is used in the fit. Also the unclustered energy scale is altered by one standard deviation and changes are propagated to $\cancel{E}_T$. These effects cause a shift of the collinear mass distribution and the resulting distribution is used in the fit. Systematic uncertainties which affect the Higgs-boson processes are discussed in [5]. Table 8.11 lists the normalisation uncertainties of the background processes.

Table 8.11.: The systematic uncertainties in the expected event yield in percentage. All uncertainties are treated as correlated between the categories, except when two values are quoted, in which case the number denoted by an asterisk is treated as uncorrelated between categories [5].

Systematic uncertainty	0-jet	1-jet	2-jet
Muon trigger/ID/isolation	2	2	2
Electron trigger/ID/isolation	3	3	3
$Z \rightarrow \tau\tau$ background	$3 \oplus 5^*$	$3 \oplus 5^*$	$3 \oplus 10^*$
$Z \rightarrow ee, \mu\mu$ background	30	30	30
Misidentified lepton background	40	40	40
Pileup	2	2	2
WW, WZ, ZZ+jets background	15	15	15
$t\bar{t}$ background	10	10	$10 \oplus 10^*$
tW background	25	25	25
b-tagging veto	3	3	3
Luminosity	2.6	2.6	2.6

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

Table 8.12.: Event yields in the signal region, $100 \text{ GeV} < M_{\text{col}} < 150 \text{ GeV}$, after fitting signal and background contributions to collision data. The expected signal yield is estimated assuming the SM Higgs boson production cross-section and for a branching fraction of $B(H \rightarrow e\tau) = 0.69\%$. Simulated events were normalised to the integrated luminosity of 19.7 fb^{-1} [5].

Jet category:	0-jet	1-jet	2-jet
Misidentified leptons	85.2 ± 5.9	38.1 ± 3.9	2.1 ± 0.7
$Z \rightarrow ee, \mu\mu$	2.3 ± 0.6	5.4 ± 0.5	-
$Z \rightarrow \tau\tau$	84.7 ± 2.1	113.3 ± 4.2	8.5 ± 0.6
$t\bar{t}, tW$	13.8 ± 0.3	69.4 ± 2.3	12.7 ± 0.8
WW, WZ, ZZ	83.0 ± 2.7	51.7 ± 2.0	3.6 ± 0.4
$W\gamma(^*)$	2.2 ± 1.0	1.2 ± 0.6	-
SM $H \rightarrow \tau\tau$	2.3 ± 0.3	3.6 ± 0.4	1.1 ± 0.2
Sum of background	273 ± 6.1	282.0 ± 6.0	28.1 ± 1.3
Observed	286	268	33
LFV $H \rightarrow e\tau$ signal	23.1 ± 1.6	16.0 ± 1.2	5.9 ± 1.0

Results The collinear mass distributions of the fitted signal and background contributions are shown in Figure 8.31. Corresponding event yields in the signal region are shown in Table 8.12. There is no evidence of a signal. The expected and observed 95% CL upper limits on $B(H \rightarrow e\tau)$ are listed in Table 8.13 assuming the Higgs-boson mass of $M_H = 125 \text{ GeV}$, and the SM production cross sections. An observed (expected) upper limit on $B(H \rightarrow e\tau) < 0.74\%$ (0.85% expected) can be set at 95% CL for the combination of all jet categories.

At $\sqrt{s} = 8 \text{ TeV}$, a search for LFV decays of the Higgs boson to $e\tau$ was also done in the $e\tau_h$ channel [5]. Figure 8.32 shows the expected and observed 95% CL upper limits on $B(H \rightarrow e\tau)$ in the $e\tau$ and $e\tau_h$ channels, and their combination. The combined observed (expected) upper limit on $B(H \rightarrow e\tau)$ is 0.69% (0.75%) at 95% CL [5].

Table 8.13.: The expected and observed upper limits at 95% on the branching fraction $B(H \rightarrow e\tau)$ for the different jet categories of the $e\tau_\mu$ channel [5]. The combination of the jet categories of the $e\tau_\mu$ channel is also shown. The asymmetric one standard-deviation uncertainties around the expected limits are shown in parentheses.

Jet category	0-jet	1-jet	2-jet	combination
Expected limit at 95% CL (%)	$< 1.63 \left({}^{+0.66}_{-0.44} \right)$	$< 1.54 \left({}^{+0.71}_{-0.47} \right)$	$< 1.59 \left({}^{+0.93}_{-0.55} \right)$	$< 0.85 \left({}^{+0.37}_{-0.25} \right)$
Observed limits at 95% CL (%)	< 1.83	< 0.94	< 1.49	< 0.74

8.10. Search for lepton-flavour violating Higgs-boson decays at $\sqrt{s} = 8$ TeV

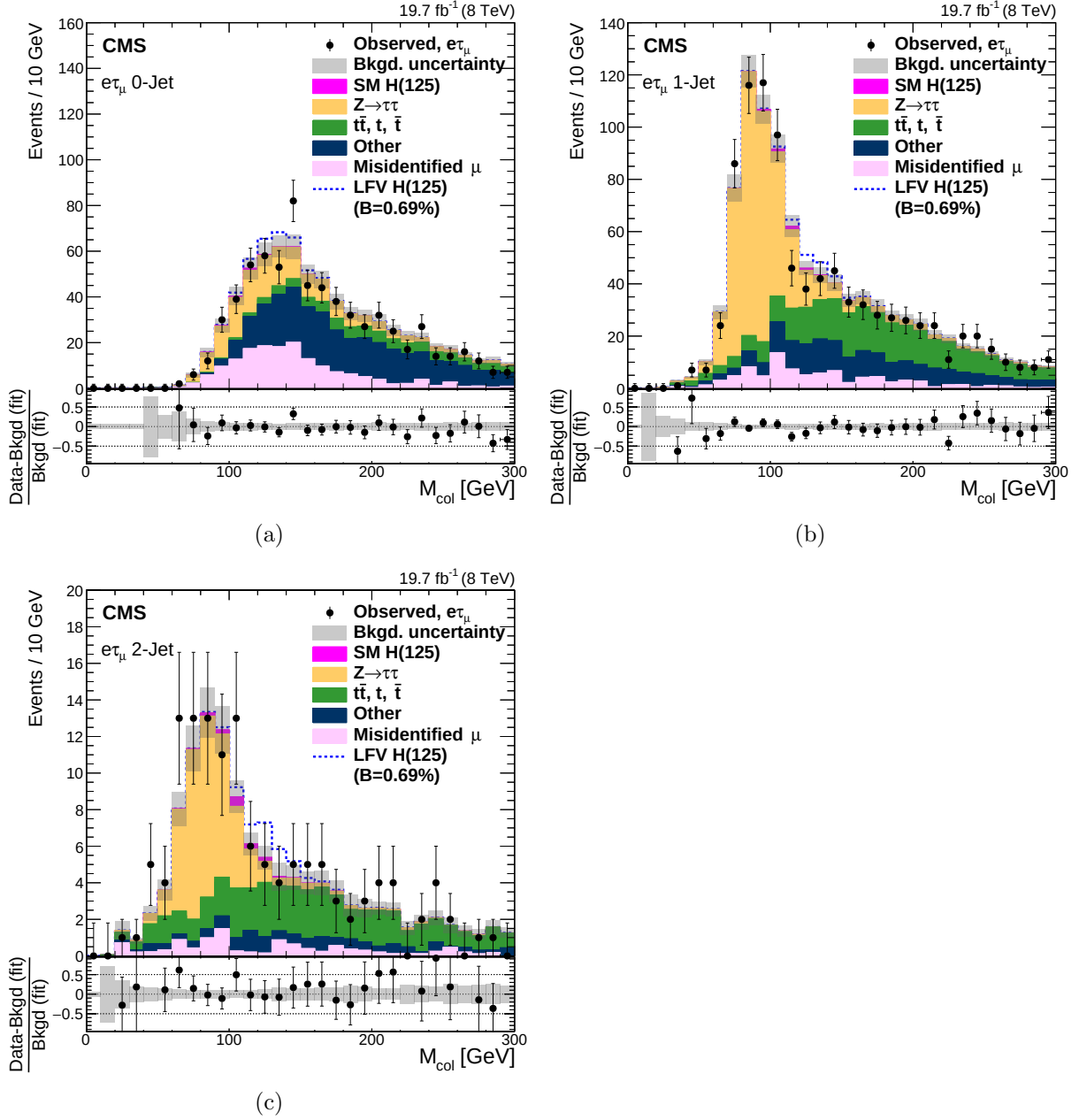


Figure 8.31.: Collinear mass distributions after event selection for the 0-jet category (a), 1-jet category (b), and 2-jet category (c). Observed distributions are compared with the background expectation after the fit. The simulated distributions for the signal are shown for the branching fraction $B(H \rightarrow e\tau) = 0.69\%$. Collision data was collected in proton-proton collisions at $\sqrt{s} = 8$ TeV [5].

Constraints on lepton flavour violating couplings The 95% CL upper limit on the branching fraction $B(H \rightarrow e\tau)$ can be interpreted in terms of the LFV Yukawa couplings $|Y_{e\tau}|$, and $|Y_{\tau e}|$ by assuming the diagonal couplings according to their SM values and negligible contributions of the other LFV Yukawa couplings to the total decay-width of the Higgs boson [5, 6]. $Y_{e\tau}$ is the coupling of the Higgs boson to a left-handed positron and a right-handed tau, while $Y_{\tau e}$ is the coupling of the Higgs boson to a left-handed anti-tau and a right-handed electron. The decay width $\Gamma(H \rightarrow e\tau)$ in terms of the Yukawa couplings is given by [6]:

$$\Gamma(H \rightarrow e\tau) = \frac{M_H}{8\pi} (|Y_{\tau e}|^2 + |Y_{e\tau}|^2) \quad (8.11)$$

and the branching fraction by:

$$B(H \rightarrow e\tau) = \frac{\Gamma(H \rightarrow e\tau)}{\Gamma(H \rightarrow e\tau) + \Gamma_{\text{SM}}}, \quad (8.12)$$

with the SM Higgs boson decay width of $\Gamma_{\text{SM}} = 4.1$ MeV for a 125 GeV Higgs boson [144]. The observed 95% upper limit on $B(H \rightarrow e\tau)$ of the combined of $e\tau_\mu$ with $e\tau_h$ constrains the corresponding Yukawa couplings to $\sqrt{|Y_{\tau e}|^2 + |Y_{e\tau}|^2} < 2.4 \times 10^{-3}$ [5]. Figure 8.33 compares this result to previous indirect measurements. This constraint of the direct search for LFV decays of the Higgs boson to $e\tau$ improves the constraints on the corresponding Yukawa couplings by one order of magnitude with respect to the constraints from indirect measurements.

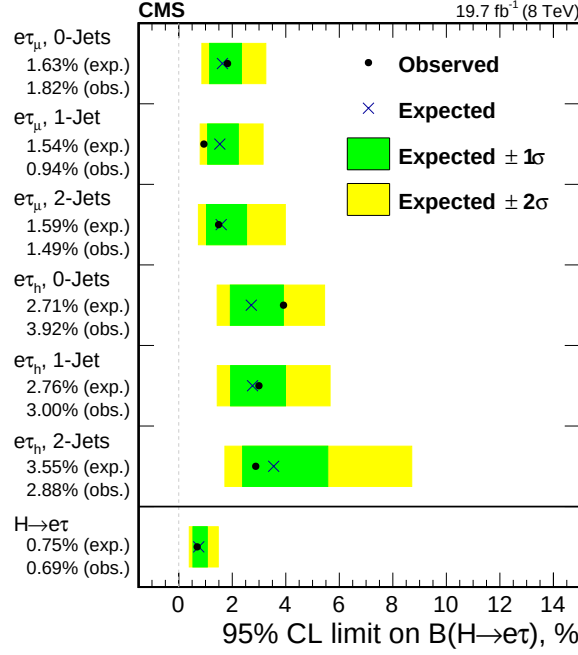


Figure 8.32.: 95% CL upper limits on the branching fraction of the LFV decay of the Higgs boson to $e\tau$. The limits by category are for a Higgs boson with $M_H = 125$ GeV assuming SM production cross sections [5].

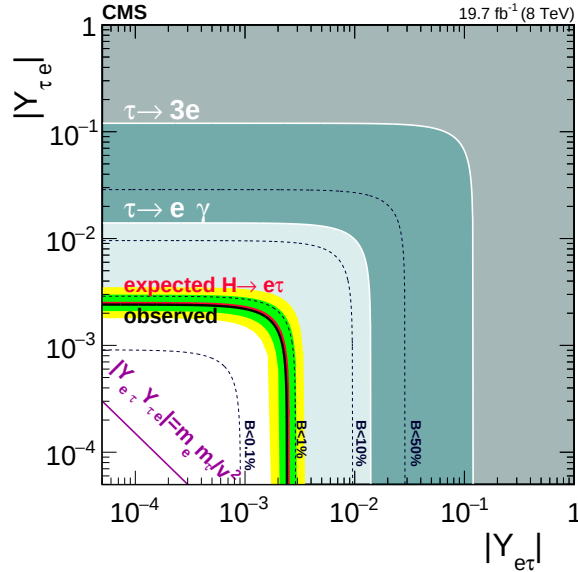


Figure 8.33.: Constraints on the lepton-flavour violating Yukawa couplings $|Y_{e\tau}|$ and $|Y_{\tau e}|$ [5]. These constraints are derived assuming the SM values for Y_{ee} , $Y_{\mu\mu}$, and $Y_{\tau\tau}$. The expected (red solid line) and the observed (black solid line) limits are derived from the 95% upper limits on the branching fraction $B(H \rightarrow e\tau)$ from the search in proton-proton collisions at $\sqrt{s} = 8$ TeV. The green (yellow) band indicates the range that is expected to contain 68% (95%) of all limit observations. The shaded regions are derived constraints from searches for $\tau \rightarrow 3e$ (grey), $\tau \rightarrow e\gamma$ (dark green), and $H \rightarrow e\tau$ (light blue). The purple diagonal line is the theoretical naturalness limit for $|Y_{e\tau} Y_{\tau e}| = \frac{m_e m_\tau}{v^2}$ [6].

8.11. Optimisation of the event selection

A new optimisation of the event selection is done due to the different centre-of-mass energy of $\sqrt{s} = 13$ TeV for the proton-proton collision and to account for the smaller statistic of the collision data with an integrated luminosity of 2.3 fb^{-1} .

First, the background modelling of the optimisation is described, followed by a description of a new optimisation method which was designed to minimise the expected 95% CL upper limit on the branching fraction. Then, the results of the optimisation using different sets of discriminating variables is described. The old optimisation method which maximises the significance $S/\sqrt{S+B}$ is redone for $\sqrt{s} = 13$ TeV. Then, the discriminating variables of the $\sqrt{s} = 8$ TeV analysis are optimised for minimising the expected 95% CL upper limit on the branching fraction. Finally, the same optimisation method is done including new discriminating variables.

8.11.1. Background modelling

Other MC samples are used in the optimisation than in the final analysis, where possible, to avoid biases in the expected upper limit coming from fluctuations in the MC samples. Drell-Yan events were simulated using MadGraph for the matrix-element calculation and PYTHIA 8 for simulating the parton shower. There was no other MC sample for the $W \rightarrow l\nu$ process with the same statistic as the sample simulated with MadGraph and PYTHIA 8. The expected event yield in the misidentified-lepton background control region is also much better than for the sample simulated using aMC@NLO for matrix element generation. As the misidentified-lepton background has the highest systematic uncertainty and is difficult to control, the aim of the event-selection optimisation is to reduce this background to a negligible amount. Thus, a bias to the expected upper limit from using the same MC sample in the optimisation and the limit calculation is expected to be small. For the $t\bar{t}$ process, events were generated which are very similar to the MC sample generated with POWHEG for the matrix element and PYTHIA 8 for the parton-shower generation. The hadron decays of the simulated $t\bar{t}$ events, which are used in the optimisation, were generated with EvtGen. Finally, contributions from the WW process were simulated with PYTHIA 8. Other background processes contribute only to a small account. To avoid biases from fluctuations in the MC samples to the event selection, further background processes, with small contributions, are not taken into account in the optimisation of the event selection. A description of the MC samples used in the optimisation is given in Table B.1 in Appendix B.

8.11.2. Optimisation method

The optimisation method used in the $e\tau_\mu$ channel of the $\sqrt{s} = 8$ TeV analysis is described in Section 8.10. This method was designed to maximise the significance $S/\sqrt{S+B}$, with S being the signal event-yield, and B the event yield of the background contributions. For the search in the $e\tau_\mu$ channel at $\sqrt{s} = 13$ TeV, the method was further developed to maximise the event selection for the expected 95% upper limit on the branching fraction of $B(H \rightarrow e\tau)$. A 95% CL upper limit corresponds to a mean of the expected upper-limit Gaussian distribution, where 95% of the distribution area would be within 1.96 ([1], Statistics, Table 39.1) standard deviations of this Gaussian. The expected upper limit on the branching fraction can be calculated by estimating how many signal events are needed to get a significance of 1.96. This number of signal events can be compared to the estimated number of signal events assuming a certain branching fraction. As a profile likelihood fit is used for estimating the expected upper limit, which takes the systematic uncertainties of the background contributions as nuisance parameters into account, a significance definition should be used in the optimisation, which also takes the systematic uncertainties of the background contributions into account. This was not the case for the significance definition used in the optimisation of the $\sqrt{s} = 8$ TeV analysis, Equation 8.10. Thus, a new significance definition is used for the $\sqrt{s} = 13$ TeV analysis. The significance Z is defined as:

$$Z = \frac{S}{\sqrt{S+B + \sum_{i=0}^b (\Delta B_i)^2}}, \quad (8.13)$$

where S and B are the number of expected signal and total background events in the mass window $100 \text{ GeV} < M_{\text{col}} < 150 \text{ GeV}$, respectively, b is the number of backgrounds, and ΔB_i is the total absolute uncertainty on the background i . Then, the expected 95% CL upper limit on the branching fraction $L_{B(H \rightarrow e\tau_\mu)}$ is given by:

$$L_{B(H \rightarrow e\tau_\mu)}[\%] = \frac{\frac{Z^2}{2} + Z \cdot \sqrt{\frac{Z^2}{4} + B + \sum_{i=0}^b (\Delta B_i)^2}}{S_{B(H \rightarrow e\tau_\mu)=1\%}}, \quad (8.14)$$

with $S_{B(H \rightarrow e\tau_\mu)=1\%}$ being the expected event yield of the signal process in the signal region for a branching-fraction of 1%. The derivation of this equation is in Appendix G. Since the systematic uncertainties of the background processes are included in the profile likelihood fit as nuisance parameters and are fitted to the observed data, the fit should have a small χ^2 for fitting the contributions of the background processes to the collision data. A

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

normalised χ^2 is be defined by:

$$\chi_{\text{normalised}}^2 = \frac{1}{B} \sum_{j=0}^{\text{bins}} \frac{\sum_{i=0}^b (\Delta B_{i,j})^2}{\sum_{i=0}^b B_{i,j}}, \quad (8.15)$$

where B is the total number of background events, $B_{i,j}$ the number of background events for background i in bin j and $\Delta B_{i,j}$ the corresponding uncertainty on the background expectation. Only bins outside the signal region mass-window are considered for the $\chi_{\text{normalised}}^2$.

Before the optimisation starts, a set of discriminating variables is defined including a step size for the following scan. First, each variable is scanned in its defined binning. There are two scans per variable. One scan which starts at the minimum value and ends at the maximal value, while the second scan goes in the other direction. At each point of the scan, a requirement on the scanning variable larger/smaller than this value is applied. The expected 95% CL upper limit is calculated and the next point in the scan is evaluated. Figure 8.34 gives two examples for such scans. After scanning the variable, the minimum value of the expected 95% CL upper limit is compared to the minimum values of the scans of the other discriminating variables. The criterion on a variable is applied, which has the smallest value for the expected 95% CL upper limit. If there are two criteria which have the same smallest value, the one with the smaller normalised χ^2 is applied. Then, all discriminating variables are scanned again. The optimisation stops, when the expected 95% CL upper limit cannot be improved, or when a requirement on the maximal number of selection criteria is reached. An improvement of this optimisation method with respect to the previous one is that the optimisation does not depend on the expected signal yield for a certain branching fraction $B(H \rightarrow e\tau_\mu)$ unlike the old method optimised for significance.

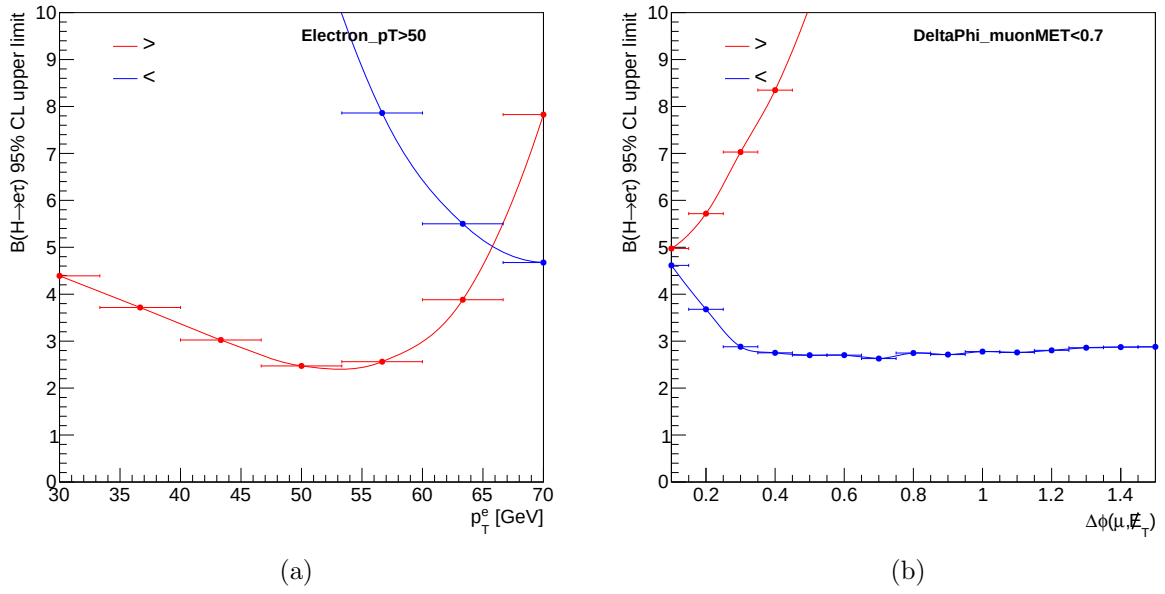


Figure 8.34.: Example variable scans for the optimisation which minimises the expected 95% CL upper limit on the branching fraction. This limit is shown as function of criteria applied on the discriminating variables p_T^e (a), and $\Delta\phi(\mu, \cancel{E}_T)$ (b). The red solid line indicates the scan which requires larger values as selection criterion, while the blue solid line indicates the other scan, which requires smaller values as selection criterion.

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

Table 8.14.: Ranking of the optimisation for maximising the significance $S/\sqrt{S+B}$ in the 0-jet category.

Variable	$S/\sqrt{S+B}$	S/B	S	B
1) $\Delta\phi(\mu, \cancel{E}_T) < 0.7$	1.33 ± 0.07	0.146 ± 0.017	13.83 ± 0.24	95 ± 11
2) $p_T^e > 50$ GeV	1.74 ± 0.1	0.378 ± 0.059	11.01 ± 0.22	29 ± 5
3) $\Delta\phi(e, \mu) > 2.5$	1.79 ± 0.11	0.444 ± 0.077	10.37 ± 0.21	23 ± 4

8.11.3. Optimisation for maximising the significance

The optimisation for maximising the old significance definition $S/\sqrt{S+B}$ used at $\sqrt{s} = 8$ TeV with the set of $\sqrt{s} = 8$ TeV discriminating variables is done to check if another event selection would have a larger expected significance for the $\sqrt{s} = 13$ TeV collision data. This is necessary for being able to compare the old selection optimisation method with the new one. A branching fraction of $B(H \rightarrow e\tau_\mu) = 1\%$ was used for the optimisation, which is in the order of the expected 95% CL upper limit of the $\sqrt{s} = 8$ TeV search.

For each jet category, the event selection is optimised. A selection of $p_T^\mu > 15$ GeV is applied to reduce the contributions from misidentified leptons, Figure C.3(b) in Appendix C. Furthermore, the electron p_T is required to be larger than 40 GeV, which is the criterion used in the loose selection of the $\sqrt{s} = 8$ TeV analysis. The following variables are scanned: p_T^e , p_T^μ , $m_T(\mu, \cancel{E}_T)$, $\Delta\phi(\mu, \cancel{E}_T)$, and $\Delta\phi(e, \mu)$. Further discriminating variables are included for the 2-jet category: $|\Delta\eta|$, and M_{jj} . Several scans of the discriminating variables were done, also without the $p_T^\mu > 15$ GeV criterion. Only discriminating variables with good results are used in the final scan. As $\Delta\phi(\mu, \cancel{E}_T)$ and $m_T(\mu, \cancel{E}_T)$ are correlated and the additional discriminating power of using both variables was found to be small, only one of them is used in the final scan. Individual scans were checked and only criteria are used to define the event selection, which have a significant impact and which do not come from fluctuations in the background contributions.

Table 8.14 shows the ranking result of the 0-jet category. The expected significance, the signal to background ratio, and the signal and background yields are listed. Figure 8.35 shows the corresponding scans. The criterion on $\Delta\phi(e, \mu)$ is not used to define the event selection, since the maximum of the significance distribution looks like coming from a fluctuation. For the 2-jet category, the statistic of the background samples was too small for using this method. Selection criteria were selected which cut the whole background away and significantly reducing the signal contributions at the same time. Thus, the optimisation was used only to select a set of discriminating variables and the selection criteria were chosen according to the distribution of the discriminating variable.

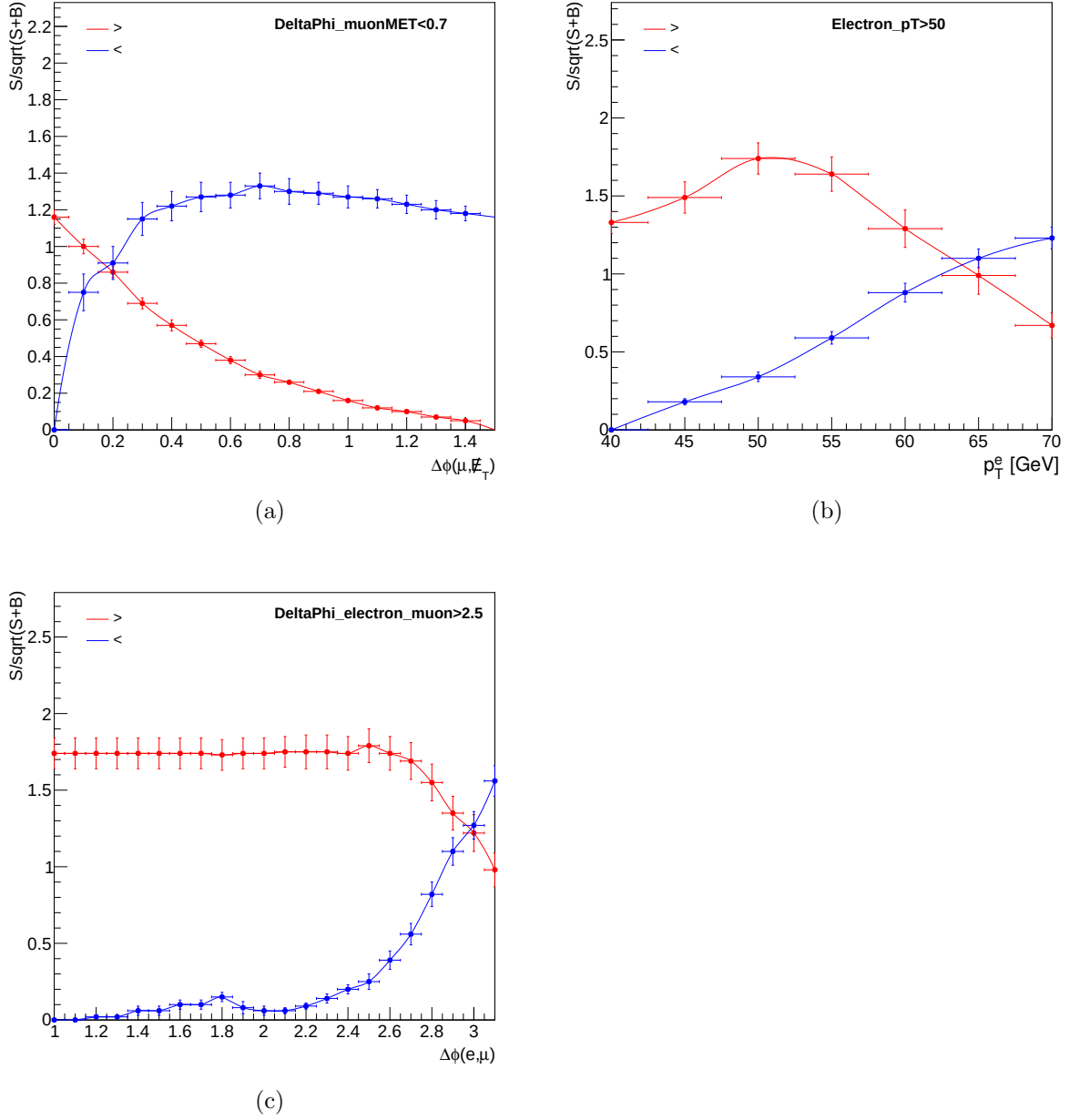


Figure 8.35.: Optimisation scans for maximising the significance in the 0-jet category. The significance $S/\sqrt{S+B}$ is shown as function of selection criteria applied on $\Delta\phi(\mu, \cancel{E}_T)$ in the first iteration (a), p_T^e in the second iteration (b), and $\Delta\phi(e, \mu)$ in the third iteration (c).

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

Table 8.15.: Event selection optimised for maximising significance using the discriminating variables of the $\sqrt{s} = 8$ TeV analysis.

Variable	0-jet	1-jet	2-jet
p_T^e	> 50 GeV	> 40 GeV	> 40 GeV
p_T^μ	> 15 GeV	> 15 GeV	> 15 GeV
$m_T(\mu, \cancel{E}_T)$	-	< 25 GeV	< 30 GeV
$\Delta\phi(\mu, \cancel{E}_T)$	< 0.7	-	-
$ \Delta\eta $	-	-	> 2

Table 8.15 lists the final event selection using the old optimisation method and the discriminating variables from the $\sqrt{s} = 8$ TeV analysis. Similar criteria as for the $\sqrt{s} = 8$ TeV optimisation are the result of the new optimisation. The requirements on $\Delta\phi(\mu, \cancel{E}_T)$ and $m_T(\mu, \cancel{E}_T)$ are tighter for the new selection, while requirements on $\Delta\phi(\mu, \cancel{E}_T)$ and $\Delta\phi(e, \mu)$ are removed for the 1-jet category. Furthermore, the VBF requirements are looser than for $\sqrt{s} = 8$ TeV to account for the smaller statistic of the 2015 collision data with respect to the $\sqrt{s} = 8$ TeV one.

8.11.4. Optimisation for minimising the expected 95% CL upper limit

The event selection was optimised again for the $\sqrt{s} = 8$ TeV variable set. For this optimisation, the expected 95% CL upper limit on the branching fraction $B(H \rightarrow e\tau_\mu)$ was minimised. A selection of $p_T^\mu > 15$ GeV is applied in the final optimisation to reduce the contributions from misidentified leptons. Electrons are required to have $p_T^e > 30$ GeV. Final scans of the 0-jet category are shown in Figure 8.36. Same criteria as in the previous optimisation are found.

Table 8.16 shows the resulting event selection. Similar selection criteria were found as for the optimisation maximising the old significance definition. Criterion on $\Delta\phi(e, \mu)$ are introduced in the 0-jet and 1-jet category, which are tighter than the criterion used in the 1-jet category in the $\sqrt{s} = 8$ TeV analysis. Additionally, the criterion on $m_T(\mu, \cancel{E}_T)$ of the 2-jet category is tighter than for the optimisation of the significance, where the criteria had to be chosen from the distributions of the discriminating variables for the 2-jet category. However, the other criteria of the 2-jet category have not changed.

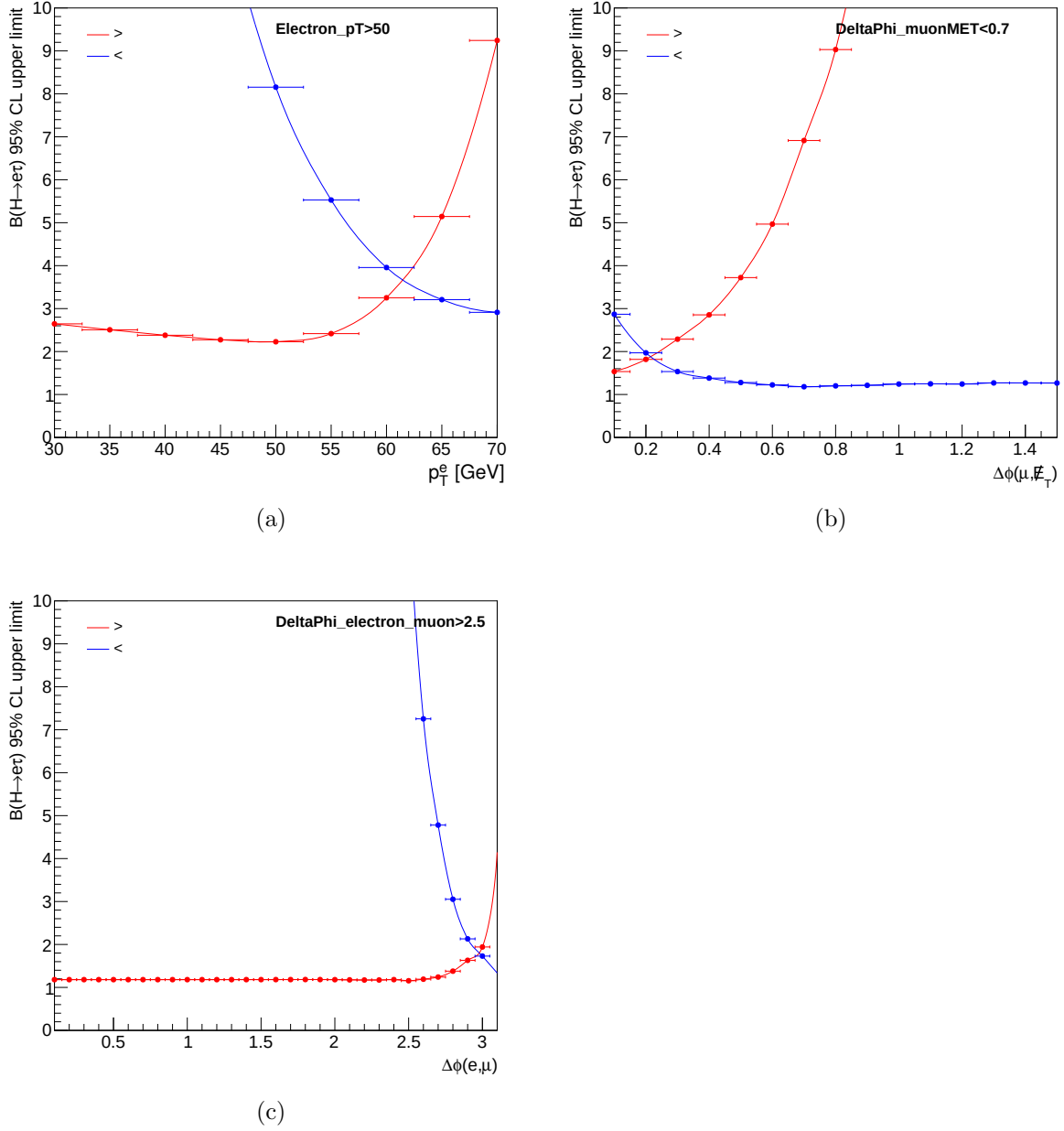


Figure 8.36.: Optimisation scans for minimising the expected 95% CL upper limit on the branching fraction of $B(H \rightarrow e\tau)$ in the 0-jet category. The expected 95% CL upper limit on the branching fraction is shown as function of selection criteria applied on p_T^e in the first iteration (a), $\Delta\phi(\mu, \cancel{E}_T)$ in the second iteration (b), and $\Delta\phi(e, \mu)$ in the third iteration (c).

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

Table 8.16.: Event selection optimised for minimising the expected 95% upper limit on the branching fraction $B(H \rightarrow e\tau_\mu)$ using the discriminating variables of the $\sqrt{s} = 8$ TeV analysis.

Variable	0-jet	1-jet	2-jet
p_T^e	$50 > \text{GeV}$	$> 40 \text{ GeV}$	$> 30 \text{ GeV}$
p_T^μ	$> 15 \text{ GeV}$	$> 15 \text{ GeV}$	$> 15 \text{ GeV}$
$m_T(\mu, \cancel{E}_T)$	-	$< 25 \text{ GeV}$	$< 20 \text{ GeV}$
$\Delta\phi(\mu, \cancel{E}_T)$	< 0.7	-	-
$\Delta\phi(e, \mu)$	> 2.5	> 0.9	-
$ \Delta\eta $	-	-	> 2

8.11.5. Optimisation for expected limit using new discriminating variables

New discriminating variables are included in the optimisation. These variables are: $\cancel{E}_{T,\text{proj}}$, $x_{\tau_\mu}^{\text{vis}}$, A_ϕ , and $x_{\cancel{E}_{T,\text{proj}}}$. A positive value is required for $\cancel{E}_{T,\text{proj}}$. This criterion is removed in the final selection, when a stronger criterion is applied on $\Delta\phi(\mu, \cancel{E}_T)^3$. Table 8.17 shows the final results of this optimisation. The criterion on the electron p_T of the 1-jet category and the criterion on $m_T(\mu, \cancel{E}_T)$ of the 2-jet category are looser when including the new variables. All other criteria are the same for the $\sqrt{s} = 8$ TeV discriminating variables. An additional criterion on the visible fraction $x_{\tau_\mu}^{\text{vis}}$ can improve the expected limit of the 0-jet and 1-jet categories. The asymmetry variable A_ϕ can improve the limit in the 1-jet and 2-jet categories. For the 2-jet category, a criterion for $x_{\cancel{E}_{T,\text{proj}}}$ is also applied.

Table 8.17.: Event selection optimised for minimising the expected 95% CL upper limit on $B(H \rightarrow e\tau_\mu)$ using new discriminating variables.

Variable	0-jet	1-jet	2-jet
p_T^e	$> 50 \text{ GeV}$	$> 35 \text{ GeV}$	$> 30 \text{ GeV}$
p_T^μ	$> 15 \text{ GeV}$	$> 15 \text{ GeV}$	$> 15 \text{ GeV}$
$m_T(\mu, \cancel{E}_T)$	-	$< 25 \text{ GeV}$	$< 25 \text{ GeV}$
$\Delta\phi(\mu, \cancel{E}_T)$	< 0.7	-	-
$\Delta\phi(e, \mu)$	> 2.5	> 0.8	-
$ \Delta\eta $	-	-	> 2
$\cancel{E}_{T,\text{proj}}$	-	> 0	> 0
$x_{\tau_\mu}^{\text{vis}}$	< 0.8	< 0.9	-
A_ϕ	-	> 0.2	> -0.2
$x_{\cancel{E}_{T,\text{proj}}}$	-	-	> 0.2

³The sign of $\cancel{E}_{T,\text{proj}}$ is strongly correlated with $\Delta\phi(\mu, \cancel{E}_T)$. It is positive for $\Delta\phi(\mu, \cancel{E}_T) < 90^\circ \approx 1.57$. Thus, if a stronger criterion on $\Delta\phi(\mu, \cancel{E}_T)$ is applied, $\cancel{E}_{T,\text{proj}}$ is always positive at the same time.

8.12. Expected 95% CL upper limits on the branching fraction $B(H \rightarrow e\tau_\mu)$

In the following, the sensitivity of the different event selections of the previous section is studied. This is done by estimating the expected 95% CL upper limits on the branching fraction $B(H \rightarrow e\tau_\mu)$ for each event selection. Here, the expected upper limit is estimated from MC simulation only, without fitting the background expectations to collision data.

First, the results for the event selection of the $\sqrt{s} = 8$ TeV search applied to simulation for $\sqrt{s} = 13$ TeV proton-proton collisions are discussed. Then, the limits of using the event selection optimised for maximising the significance using expected background contributions for $\sqrt{s} = 13$ TeV are shown. Afterwards, the limits for the event selection which was optimised to minimise the expected limit is discussed. Finally, the results are presented for the event selection which includes new discriminating variables.

Event selection of the $\sqrt{s} = 8$ TeV analysis: The search for lepton-flavour violating decays of the Higgs boson to $e\tau$ in the $e\tau_\mu$ channel using the $\sqrt{s} = 8$ TeV proton-proton collision data is described in Section 8.10. This search uses an event selection which was optimised to maximise the significance $S/\sqrt{S+B}$. Table 8.10 summarises the event selection for each category. Figure F.1 in Appendix F shows the event yield at $\sqrt{s} = 13$ TeV after preselection as function of the discriminating variables used to define the event selection of the 0-jet category in the $\sqrt{s} = 8$ TeV analysis. The criteria on the variables are indicated as black line and an arrow indicates the considered events. Event yields of the 1-jet category are shown in Figure F.2, and Figure F.3, while Figure F.4 and Figure F.5 show the corresponding event yields for the 2-jet category. The event selection of the $\sqrt{s} = 8$ TeV analysis can also be applied at $\sqrt{s} = 13$ TeV, but the VBF requirements seem to be a bit tight.

Figure 8.37 shows the event yield as function of the collinear mass for each jet category after applying the $\sqrt{s} = 8$ TeV event selection. The statistics of the 0-jet and 1-jet category are sufficient for a template fit, while the statistic of the 2-jet category is quite low. The expected 95% CL upper limits on the branching fraction $B(H \rightarrow e\tau)$ are listed in Table 8.18. Only the expected limit of the 0-jet category is comparable to the ones at $\sqrt{s} = 8$ TeV, Table 8.13. The limit is a bit improved, while the 68% limit regions are similar. For the 1-jet category, the expected limit is worse and the 68% region is increased with respect to $\sqrt{s} = 8$ TeV. This is even worse for the 2-jet category.

Table 8.18.: Expected upper limits at 95% CL on the branching fraction $B(H \rightarrow e\tau)$ for the event selection which was optimised for $\sqrt{s} = 8$ TeV analysis. The expected upper limits were derived for $\sqrt{s} = 13$ TeV.

Category	Expected limits at 95% CL (%)
0-jet	$1.55^{+0.69}_{-0.46}$
1-jet	$3.23^{+1.48}_{-0.98}$
2-jet	$5.58^{+2.78}_{-1.78}$
combined	$1.32^{+0.58}_{-0.39}$

8.12. Expected 95% CL upper limits on the branching fraction $B(H \rightarrow e\tau_\mu)$

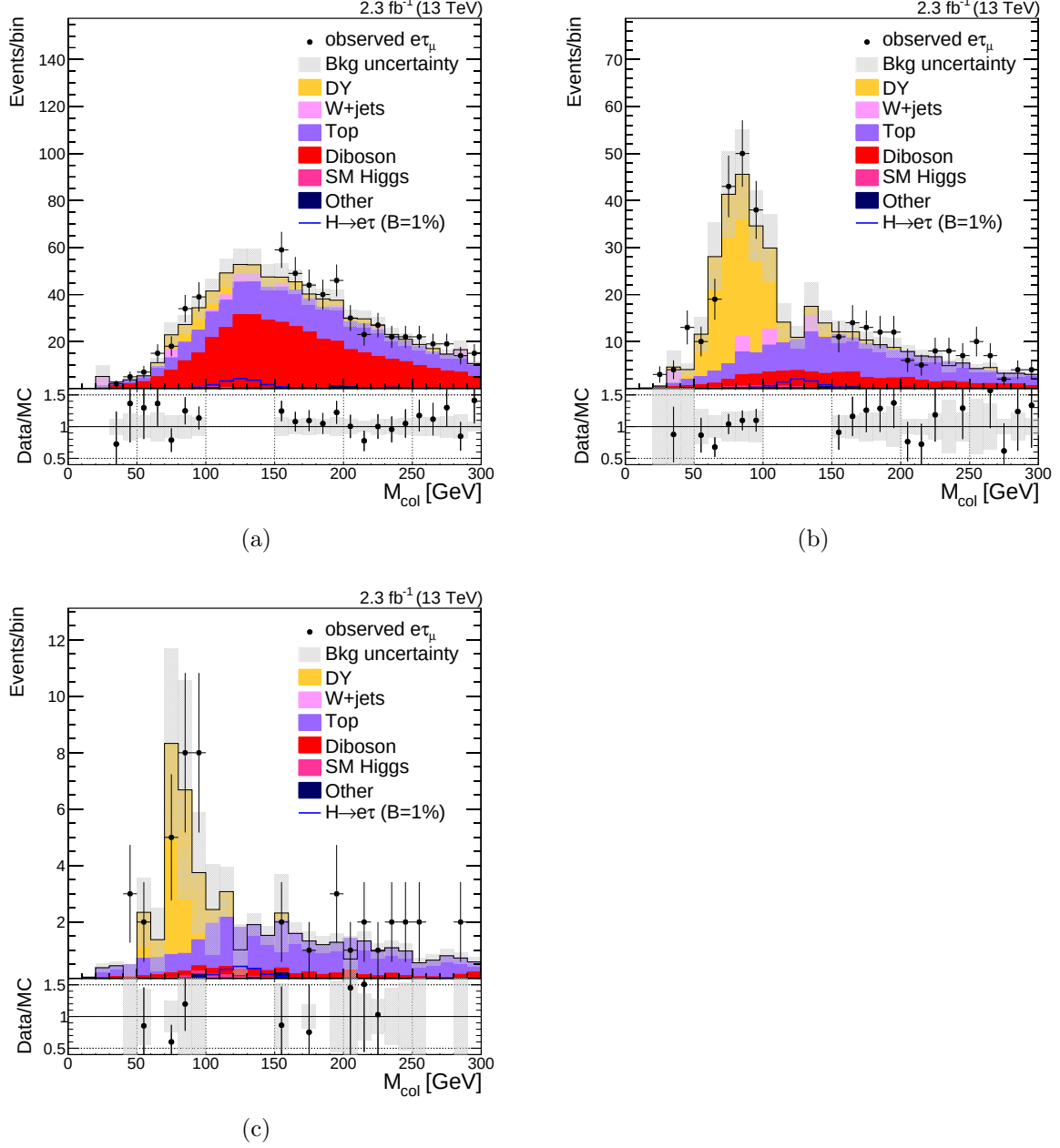


Figure 8.37.: Collinear mass of the event selection which was optimised for $\sqrt{s} = 8$ TeV. The event yield is shown as function of the collinear mass for the 0-jet category (a), 1-jet category (b), and 2-jet category (c). Signal regions are blinded and collision data were collected in proton-proton collisions at $\sqrt{s} = 13$ TeV.

Event selection optimised for significance: The event selection of the $\sqrt{s} = 8$ TeV analysis is optimised again for the $\sqrt{s} = 13$ TeV collision data. Changes to the criteria are small except for the less tight VBF criteria. Figure 8.38 shows the collinear mass for each jet category applying the event selection. Table 8.19 shows the expected 95% CL upper limits. They are comparable to the ones when applying the $\sqrt{s} = 8$ TeV event selection. The new event selection has better expected limits for the 0-jet and 2-jet category, while the one for the 1-jet category is worse. In summary, the combined expected limit is improved with respect to the expected limit when using the $\sqrt{s} = 8$ TeV event selection. However, all improvements are within 1.96 standard deviations and not very significant.

Table 8.19.: Expected 95% CL upper limits on the branching fraction $B(H \rightarrow e\tau)$ for the event selection which was optimised for maximising the significance using the discriminating variables of the $\sqrt{s} = 8$ TeV analysis. The expected upper limits were derived for $\sqrt{s} = 13$ TeV.

Category	Expected limits at 95% CL (%)
0-jet	$1.49^{+0.68}_{-0.45}$
1-jet	$3.55^{+1.58}_{-1.04}$
2-jet	$5.20^{+2.55}_{-1.64}$
combined	$1.30^{+0.58}_{-0.39}$

8.12. Expected 95% CL upper limits on the branching fraction $B(H \rightarrow e\tau_\mu)$

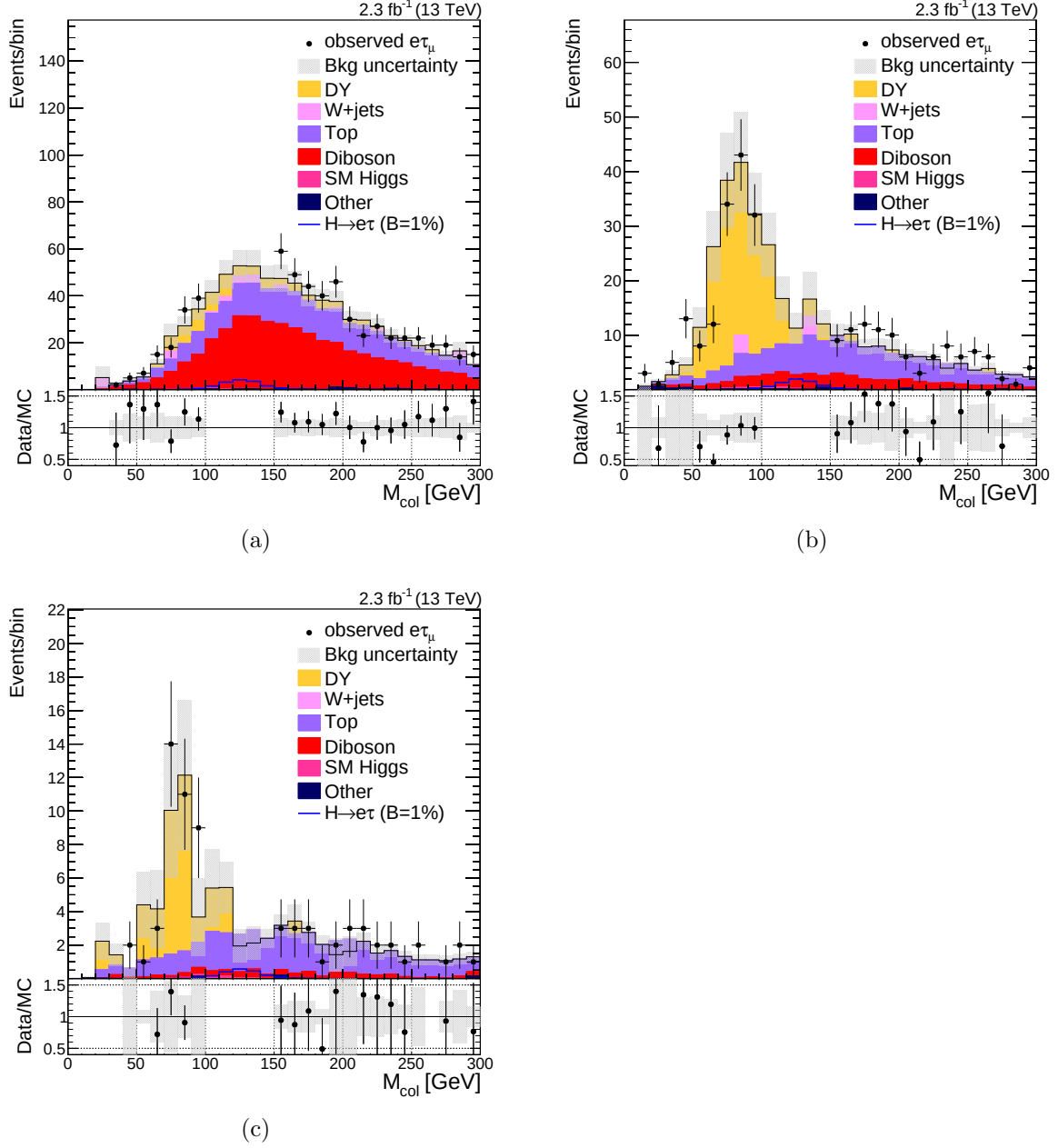


Figure 8.38.: Collinear mass of the event selection which was optimised for maximising significance using the discriminating variables of the $\sqrt{s} = 8$ TeV analysis. The event yield is shown as function of the collinear mass for the 0-jet category (a), 1-jet category (b), and 2-jet category (c). Signal regions are blinded and collision data were collected in proton-proton collisions at $\sqrt{s} = 13$ TeV.

Event selection optimised for expected limit: The event selection is also optimised for minimising the expected 95% upper limit on the branching fraction $B(H \rightarrow e\tau)$ using the discriminating variables of the $\sqrt{s} = 8$ TeV search. Figure 8.39 shows the collinear mass distribution after applying the new event selection, and Table 8.20 lists the expected 95% CL upper limits on the branching fraction, which are comparable to the limits of the previous two event selections. The mean limits on the 1-jet and 2-jet category are smaller, while the mean limit on the 0-jet category is larger. This comes from the criterion applied on $\Delta\phi(e, \mu)$. The mean of the combined expected limit is smaller than for the previous event selections. However, this improvement is not significant.

Table 8.20.: Expected upper limits at 95% CL on the branching fraction $B(H \rightarrow e\tau)$ for the event selection which was optimised for minimising the expected upper limit using the discriminating variables of the $\sqrt{s} = 8$ TeV analysis. The expected upper limits were derived for $\sqrt{s} = 13$ TeV.

Category	Expected limits at 95% CL (%)
0-jet	$1.51^{+0.69}_{-0.45}$
1-jet	$3.10^{+1.45}_{-0.95}$
2-jet	$4.80^{+2.39}_{-1.53}$
combined	$1.26^{+0.57}_{-0.38}$

8.12. Expected 95% CL upper limits on the branching fraction $B(H \rightarrow e\tau_\mu)$

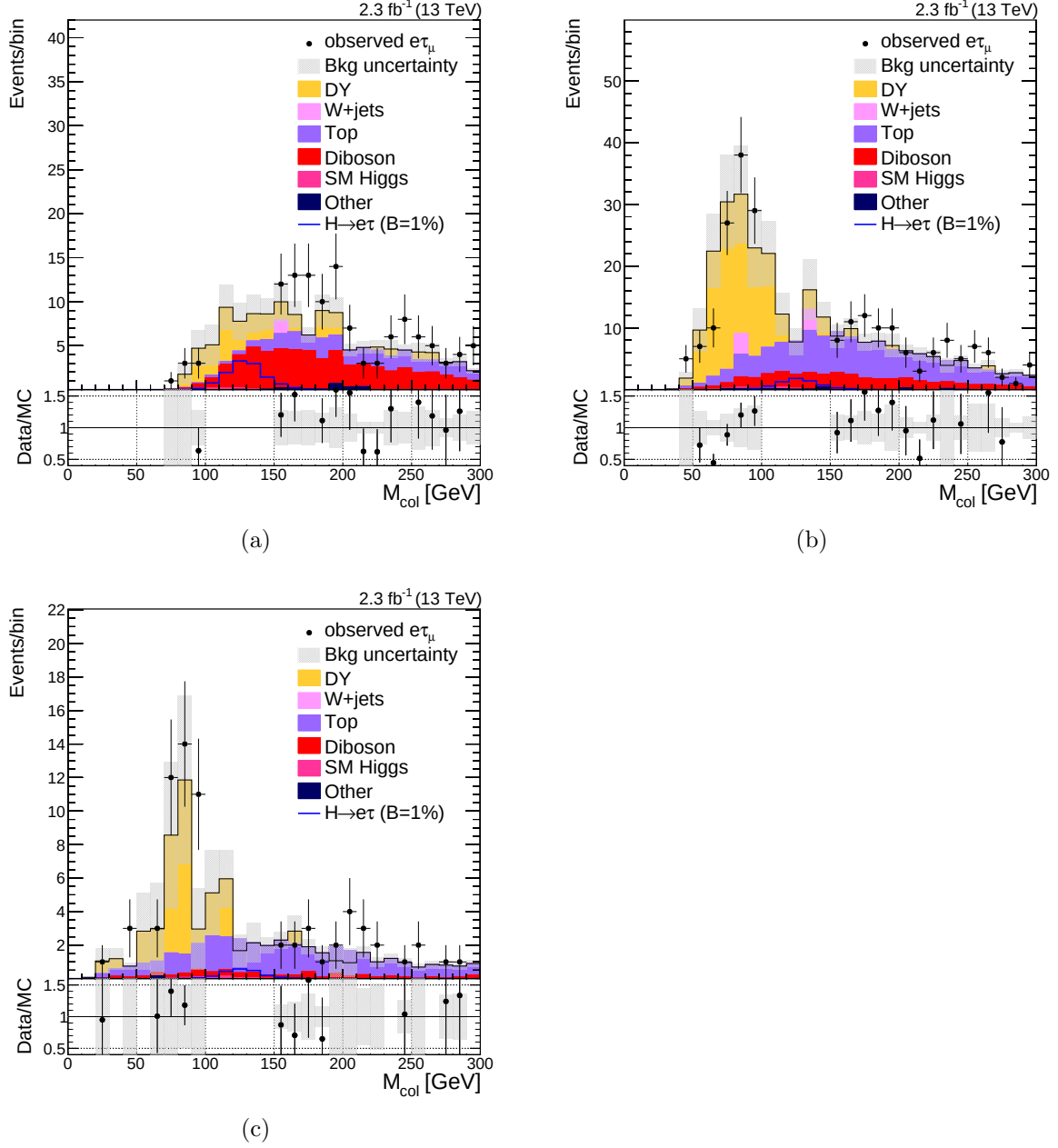


Figure 8.39.: Collinear mass of the event selection which was optimised for minimising the expected 95% CL upper limit on $B(H \rightarrow e\tau)$ using the discriminating variables of the $\sqrt{s} = 8$ TeV analysis. The event yield is shown as function of the collinear mass for the 0-jet category (a), 1-jet category (b), and 2-jet category (c). Signal regions are blinded and collision data were collected in proton-proton collisions at $\sqrt{s} = 13$ TeV.

Event selection optimised for expected limit using new discriminating variables: In the last optimisation, the event selection is optimised for minimising the expected 95% CL upper limit on the branching fraction $B(H \rightarrow e\tau)$ including new discriminating variables. The distributions of the collinear mass is shown in Figure 8.40 for each jet category after applying the event selection. Table 8.21 gives a summary of the expected 95% CL upper limits on the branching fraction. They are comparable to the limits of the other event selections. The limits of the 0-jet and 1-jet category are better than for the previous event selections. For the 2-jet category, the upper expected limit is worse than for the other event selections. The combined expected limit is better than for the previous event selections. However, the improvement is within 1.96 standard deviations.

Table 8.21.: Expected upper limits at 95% CL on the branching fraction $B(H \rightarrow e\tau)$ for the event selection which was optimised for minimising the expected upper limit on the branching fraction using new discriminating variables. The expected upper limits were derived for $\sqrt{s} = 13$ TeV.

Category	Expected limits at 95% CL (%)
0-jet	$1.46^{+0.67}_{-0.44}$
1-jet	$2.91^{+1.38}_{-0.89}$
2-jet	$5.23^{+2.40}_{-1.61}$
combined	$1.22^{+0.55}_{-0.36}$

Resulting event selection for $\sqrt{s} = 13$ TeV: The event selection including the new discriminating variables has the best expected limit. As the event selection without the new variables has a better limit and fewer selection criteria for the 2-jet category, this selection criteria are used for the 2-jet category. The criteria on the new discriminating variables are removed and the looser criterion on the transverse mass is kept.

8.12. Expected 95% CL upper limits on the branching fraction $B(H \rightarrow e\tau_\mu)$

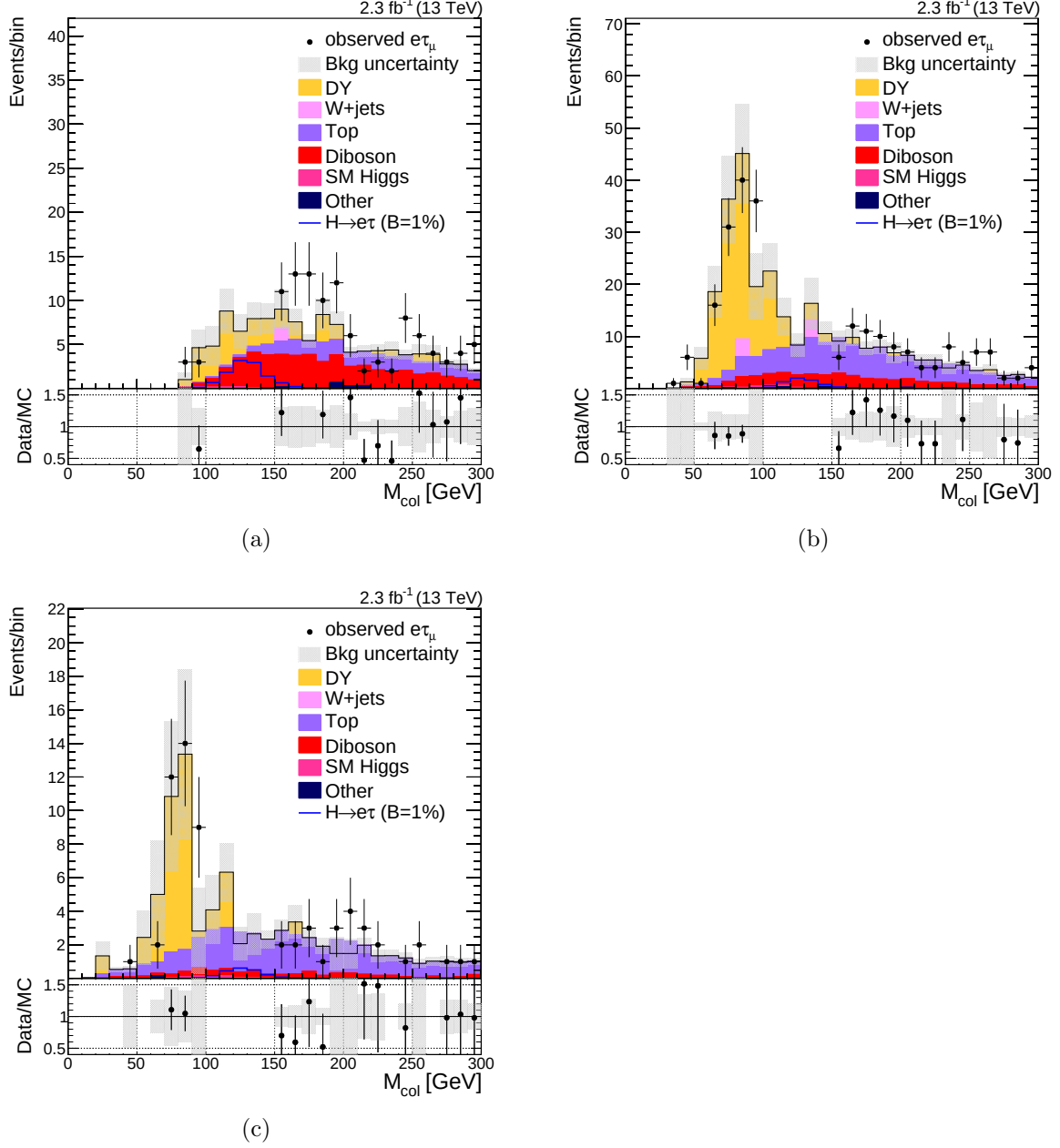


Figure 8.40.: Collinear mass of the event selection which was optimised for expected 95% CL upper limit on $B(H \rightarrow e\tau)$ using new discriminating variables. The event yield is shown as function of the collinear mass for the 0-jet category (a), 1-jet category (b), and 2-jet category (c). Signal regions are blinded and collision data were collected in proton-proton collisions at $\sqrt{s} = 13$ TeV.

8.13. Event selection

The final event selection for the 2015 run at $\sqrt{s} = 13$ TeV is summarised in Table 8.22. Figure 8.41 shows the collinear mass distributions of the estimated signal and background contributions. In the 0-jet category, the signal region is dominated by the diboson background contributions. There are also contributions from the $Z \rightarrow ll$ and the top processes. Contributions from misidentified leptons are negligible. The signal region of the 1-jet category is dominated by contributions from the top processes. There are some contributions from the diboson processes and the Drell-Yan process. For the 2-jet category, the signal region is dominated by contributions from top processes. The $Z \rightarrow ll$ process and the diboson processes also contribute.

Table 8.22.: Final event selection of the search for LFV decays of the Higgs boson to $e\tau_\mu$ at $\sqrt{s} = 13$ TeV.

Variable	0-jet	1-jet	2-jet
p_T^e	> 50 GeV	> 35 GeV	> 30 GeV
p_T^μ	> 15 GeV	> 15 GeV	> 15 GeV
$m_T(\mu, \cancel{E}_T)$	-	< 25 GeV	< 25 GeV
$\Delta\phi(\mu, \cancel{E}_T)$	< 0.7	-	-
$\Delta\phi(e, \mu)$	> 2.5	> 0.8	-
$ \Delta\eta $	-	-	> 2
$\cancel{E}_{T,\text{proj}}$	-	> 0 GeV	-
$x_{\tau_\mu}^{\text{vis}}$	< 0.8	< 0.9	-
A_ϕ	-	> 0.2	-

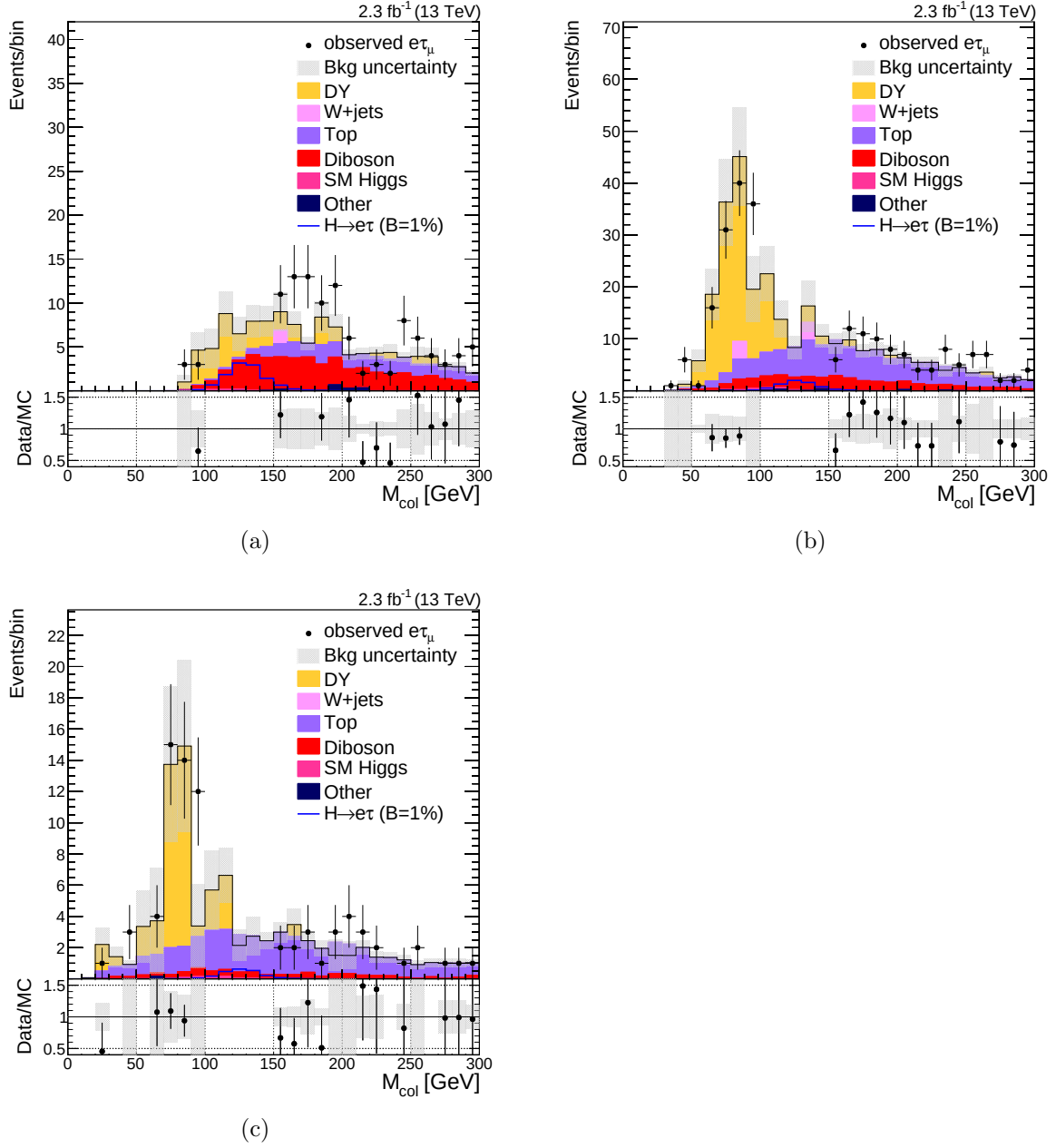


Figure 8.41.: Collinear mass distributions after event selection for the 0-jet category (a), 1-jet category (b), and 2-jet category (c). Observed distributions are compared with the background expectation. The simulated distributions for the signal are shown for the branching fraction $B(H \rightarrow e\tau)=1\%$. Collision data was collected in proton-proton collisions at $\sqrt{s} = 13$ TeV. The signal regions are blinded.

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

The expected 95% CL upper limits on $B(H \rightarrow e\tau)$ are listed in Table 8.23 assuming the Higgs-boson mass of $M_H = 125$ GeV, and the SM production cross sections. They were estimated from MC simulation, without fitting the background estimation to collision data. The combined expected upper limit on $B(H \rightarrow e\tau)$ is 1.21%. This expected upper limit is better than for the event selection optimised for $\sqrt{s} = 8$ TeV, which has a combined expected upper limit of $1.26 \left({}^{+0.57}_{-0.38} \right)$. The upper limits are slightly improved for the 0-jet and 1-jet categories with respect to the event selection optimised for $\sqrt{s} = 8$ TeV. Furthermore, the uncertainties decrease for the 1-jet category.

Table 8.23.: The expected upper limits at 95% CL on the branching fraction $B(H \rightarrow e\tau)$ for the different jet categories. The asymmetric one standard-deviation uncertainties around the expected limits are shown in parentheses.

Jet category	Expected limit at 95% CL (%)
0-jet	$1.46 \left({}^{+0.67}_{-0.44} \right)$
1-jet	$2.91 \left({}^{+1.38}_{-0.89} \right)$
2-jet	$4.73 \left({}^{+2.32}_{-1.49} \right)$
combined	$1.21 \left({}^{+0.55}_{-0.36} \right)$

8.14. Results

The event selection described in the previous section is applied in order to enhance the expected signal with respect to the SM background contributions. It was optimised to minimise the expected upper limit on the branching fraction $B(H \rightarrow e\tau)$. In the next step, the collinear mass distribution is fitted with a binned profile-likelihood fit, where the systematic uncertainties are used as nuisance parameter in the fit. A background-only fit is done, where only the background contributions are fitted to collision data and the signal contributions are assumed to be zero. Another fit is also done, where background and signal contributions are fitted to the collision data, the signal-plus-background fit. Both fits are done using all three jet categories. There are nuisance parameters, which are correlated among the categories, such as the normalisation of the W +jets background, and there are nuisance parameters, which are uncorrelated to the other jet categories, like differences in the normalisation of a certain background between the jet categories. The nuisance parameters are summarised in Section 8.8.

Figure 8.42 shows the collinear mass distributions after the background-only fit and the signal-plus-background fit. There is no evidence of lepton-flavour violating decays of the Higgs boson to $e\tau$. Fitting signal and background contributions leads to a best fit value of $B(H \rightarrow e\tau) = 0.66 \left({}^{+0.38}_{-0.57} \right) \%$, which is compatible with a branching fraction of zero. The best fit value is smaller than the observed upper limit of the $\sqrt{s} = 8$ TeV search [5]. Post-fit plots for the preselection are shown in Figure H.1, Appendix H. Pre-fit plots for both preselection and final selection are shown in Figure H.2, Appendix H, after unblinding the signal region.

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

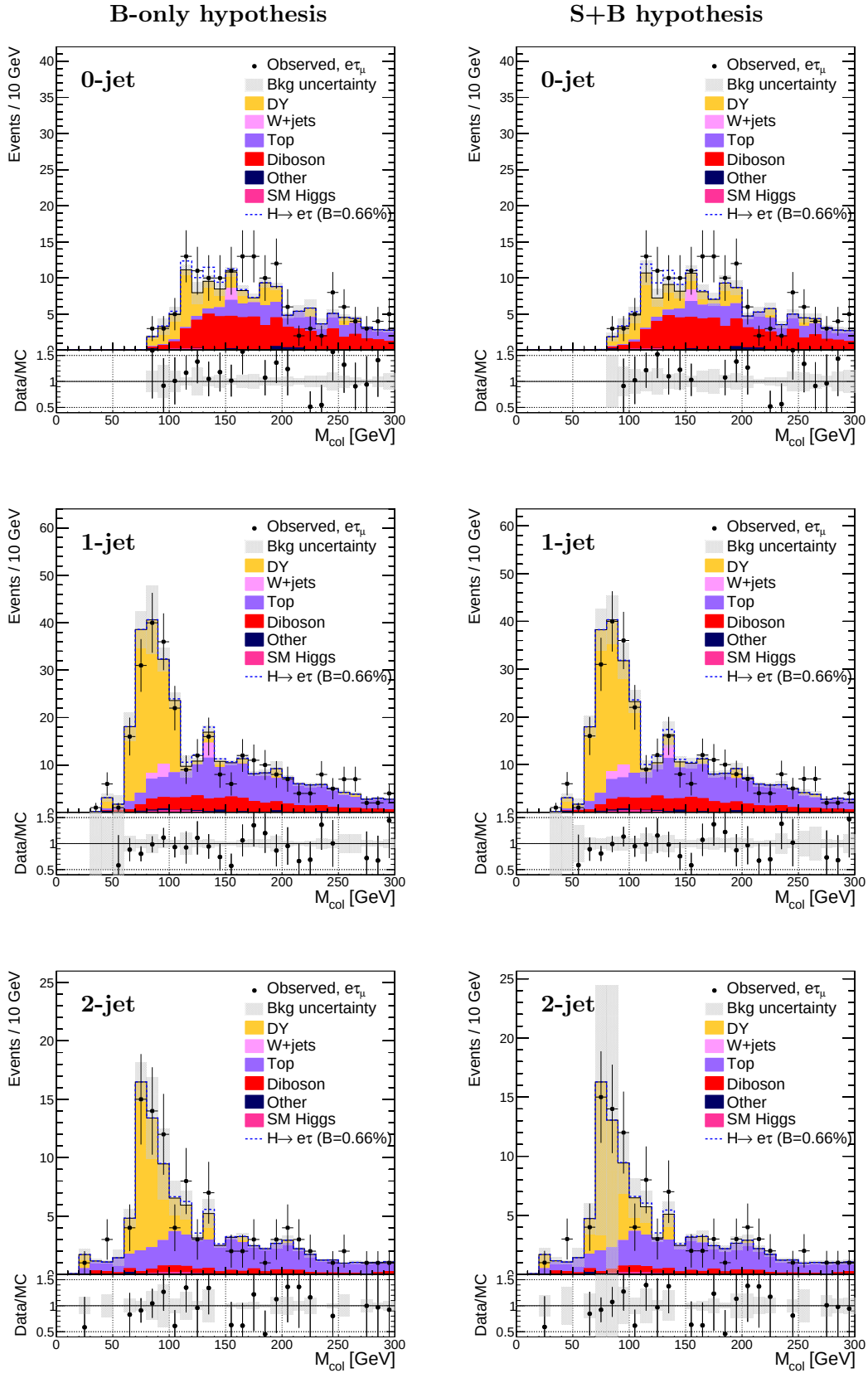


Figure 8.42.: Comparison of the observed collinear mass distribution with the background expectations after the binned likelihood fit. The signal expectation is shown for the best-fit branching fraction $B(H \rightarrow e\tau) = 0.66\%$. The left column shows the background expectation after the background-only fit, while the right column shows the background expectations after the signal-plus-background fit. In the upper, middle, and lower rows, the collinear mass is shown for the 0-jet, 1-jet, and 2-jet category, respectively. Proton-proton collision data were collected at $\sqrt{s} = 13$ TeV.

The pull distribution of the nuisance parameters, which change the normalisation of the background contributions, are shown for the preselection and the full selection in Figure 8.43. Both hypotheses, background-only and signal-plus-background, have very similar pulls in the fit, but the changes in the uncertainty on the fitted nuisance parameters behaves differently. There are also the same trends of the pull distributions in the preselection fit and the full-selection fit, except for the Z+jets background, which has a negative pull for the preselection fit and a positive pull for the full-selection fit for the background-only hypothesis. The diboson background and the unclustered energy scale are underestimated using MC simulations, while the jet energy scale is overestimated. However, these differences are in the order of one standard deviation. It is expected that the pull distributions in the jet energy scale and unclustered energy scale are anti-correlated. Furthermore, the JES is calibrated to a certain composition of jet flavours. As the response of the detector depends on the jet flavour and the jet-flavour composition of this selection may differ from the one used in calibration of the JES, differences in the JES within one standard deviation are possible. It can be seen in Figure H.2, Appendix H, that the normalisation of the diboson background is underestimated for the preselection and the 0-jet category.

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

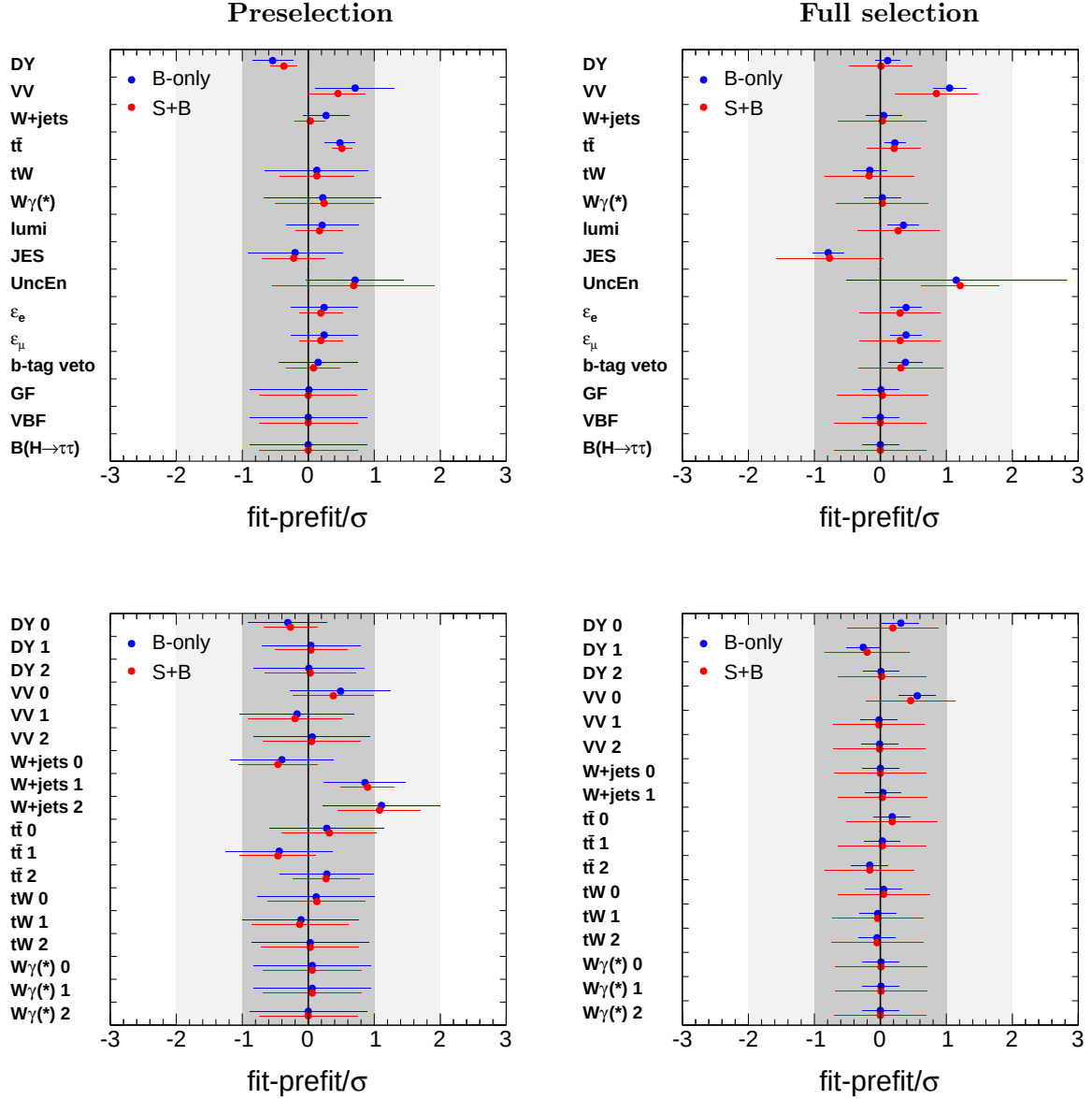


Figure 8.43.: The pull distributions are shown for the background-only (blue) and the signal-plus-background (red) hypotheses. Changes of the nuisance parameters with respect to the prefit value are shown in units of the corresponding standard deviation σ . In the fit, the uncertainty on the fitted value is also determined. Changes in the uncertainty are represented as error bars with its length smaller or larger than 1σ . The left column shows the fit pull distribution for the fits after preselection, while the right column shows the fit pull distributions after full selection. In the upper row, the nuisance parameters are shown, which affects all jet categories, while in the lower row the nuisance parameters are shown, which only affects the indicated jet category. The diboson background (WW,WZ,WW) is indicated as VV. Other nuisance parameters are systematic uncertainties on the integrated luminosity (lumi), jet energy scale (JES), unclustered energy scale (UncEn), electron or muon efficiencies (ϵ_e or ϵ_μ), misidentification efficiency of jets as b-quark jets (b-tag veto), theoretical uncertainties on gluon fusion or VBF Higgs production (GF or VBF), and the theoretical uncertainty on the branching fraction of the Higgs boson to two taus. Systematic uncertainties which are used as nuisance parameters in the fits are summarised in Section 8.8.

Table 8.24.: Event yields in the signal region, $100 \text{ GeV} < M_{\text{col}} < 150 \text{ GeV}$, after the profile likelihood fit of the background contributions only. The estimated background contributions were normalised to 2.32 fb^{-1} before the fit. Furthermore, the SM Higgs boson production cross section was assumed, taking gluon fusion and VBF Higgs-production into account. The obtained best fit branching fraction of the signal-plus-background fit, $B(H \rightarrow e\tau) = 0.66\%$, is used for estimating the expected signal yield, which is derived from MC simulation only. The uncertainty on the signal yield does only include uncertainties for the signal model.

Jet category	0-jet	1-jet	2-jet
$Z \rightarrow ll$	21.2 ± 2.2	20.8 ± 2.9	7.7 ± 2.2
Misidentified leptons	—	3.2 ± 0.8	—
$t\bar{t}$, $tW/\bar{t}W$	3.0 ± 0.3	32.4 ± 0.9	12.9 ± 0.5
WW , WZ , ZZ	17.1 ± 0.4	13.5 ± 0.6	2.2 ± 0.3
$W\gamma^{(*)}$	0.2 ± 0.0	0.9 ± 0.3	—
SM H background	0.6 ± 0.0	1.0 ± 0.0	0.5 ± 0.0
Sum of background	42.1 ± 2.1	71.8 ± 3.5	23.3 ± 2.4
Observed	49.0 ± 7.0	67.0 ± 8.2	22.0 ± 4.7
LFV $H \rightarrow e\tau$ ($B=0.66\%$)	6.6 ± 0.4	4.1 ± 0.2	1.3 ± 0.1

The event yields in the mass range $100 \text{ GeV} < M_{\text{col}} < 150 \text{ GeV}$ are given in Table 8.24 after fitting the estimated background contributions to the collision data. In Appendix H, corresponding event yields are shown for the signal-plus-background fit in Table H.1. The observed event yield matches with the fitted background yield for both fits within the uncertainties.

Expected and observed limits on the branching fraction $B(H \rightarrow e\tau)$ can be derived. Table 8.25 shows the expected and observed upper limits at 95% CL on the branching fraction $B(H \rightarrow e\tau)$ for each jet category and their combination. These limits are also summarised in Figure 8.44. The combined observed (expected) upper limit on $B(H \rightarrow e\tau)$ is 1.94% (1.31%) at 95% CL.

Table 8.25.: The expected and observed upper limits at 95% CL on the branching fraction $B(H \rightarrow e\tau)$ for a Higgs boson with $M_H = 125 \text{ GeV}$. The limits for each jet category and the combined limit are shown. Asymmetric one standard-deviation uncertainties on the expected limits are shown in parentheses.

Jet category	0-jet	1-jet	2-jet	combined
Expected limit at 95% CL (%)	$< 1.60 \left({}^{+0.73}_{-0.48} \right)$	$< 3.20 \left({}^{+1.39}_{-0.92} \right)$	$< 5.14 \left({}^{+2.44}_{-1.61} \right)$	$< 1.31 \left({}^{+0.58}_{-0.39} \right)$
Observed limit at 95% CL (%)	< 2.23	< 3.04	< 5.94	< 1.94

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

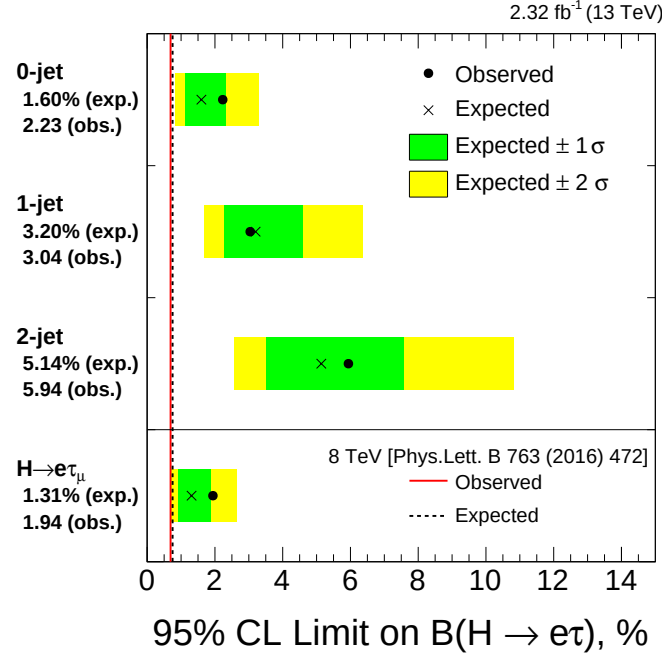


Figure 8.44.: The expected and observed upper limits at 95% CL on the branching fraction $B(H \rightarrow e\tau)$ for a Higgs boson with $M_H = 125$ GeV. The limits for each jet categories and the combined limit are shown. Asymmetric one standard-deviation uncertainties on the expected limits are shown in green, while two standard-deviation uncertainties are shown in yellow. The results of the search at $\sqrt{s} = 8$ TeV [5] are indicated as red line (observed combined limit) and black dashed line (expected combined limit).

Expected and observed limits for the analysis at $\sqrt{s} = 8$ TeV An observed (expected) upper limit on the branching fraction $B(H \rightarrow e\tau) < 0.74\%$ ($0.85^{+0.37}_{-0.25}\%$) was derived for the $e\tau_\mu$ channel using $\sqrt{s} = 8$ TeV proton-proton collision data. The $e\tau_\mu$ channel was combined with the $e\tau_h$ channel and the combined observed (expected) upper combined limit was 0.69% (0.75%) at 95% CL [5]. The three jet categories of the $e\tau_h$ channel are the most sensitive ones.

Comparison to the analysis at $\sqrt{s} = 8$ TeV in the $e\tau_\mu$ channel The statistic of the 2015 data set at $\sqrt{s} = 13$ TeV is lower than for the 2012 data set at $\sqrt{s} = 8$ TeV, which cannot be compensated by a new optimisation of the event selection. Thus, the limits of the analysis at $\sqrt{s} = 13$ TeV are weaker than the limits using the 8 TeV collision data. In addition, the uncertainties on the expected upper limit are also larger. The 13 TeV analysis has a smaller sensitivity in the 1-jet and 2-jet categories, the upper limits are two times and three times larger than for the search done at $\sqrt{s} = 8$ TeV.

A comparison between the results of both searches is done in Table H.2 and Table H.3 in Appendix H. The 2-jet category is less sensitive in the 13 TeV analysis, as the decrease of the signal is larger than for the sum of SM background contributions, mostly because

the fraction of the top background events has increased. The top events contribution has the highest systematic uncertainty among the dominant background processes, as it can be seen in Table 8.8. For the 1-jet category in the 13 TeV analysis, the increase of the fraction of the top events is even larger. It has increased by a factor of 1.8.

One reason for the decrease in sensitivity of the 1-jet and 2-jet categories could be the increase in the top background compared to the other backgrounds and the Higgs boson production. For example, the Higgs boson production via gluon fusion increased by a factor of $44.14/19.47 \approx 2$ [142,145], while the cross section of $t\bar{t}$ production increases by a factor of $772/245 \approx 3$ [129,146]. The production cross-sections of the other dominant backgrounds have also increased approximately by a factor of two with respect to $\sqrt{s} = 8$ TeV. Thus, the $t\bar{t}$ background and the reduction of top events become important. Another reason for the smaller sensitivity of the search done with the 2015 collision data set, is the reduced statistic. Table H.4 in Appendix H shows the expected number of events from SM Higgs production via gluon fusion and VBF Higgs production for the 2012 and 2015 collision-data sets. The number of expected events from $t\bar{t}$ production is also shown. It can be seen that the 2015 data set with 2.32 fb^{-1} has about one third less statistics than the 2012 data set with 19.7 fb^{-1} . This reduced statistics makes it more difficult to extract the background contributions in the template fit due to a larger sensitivity to fluctuations in the data set. Furthermore, it is more difficult to enhance the signal contributions with respect to the SM background contributions without applying an event selection without further reducing the statistics.

Expected-limit as function of the integrated luminosity at $\sqrt{s} = 13$ TeV Figure 8.44 shows the expected upper limit on the branching fraction $B(H \rightarrow e\tau)$ at 95% CL as function of the integrated luminosity at $\sqrt{s} = 13$ TeV. The upper limits were derived using only MC simulation. This is done for each jet category and for their combination. To compare the derived limits with the results of the 8 TeV analysis [5], the expected limits for the $e\tau_\mu$ channel and for the combination with the $e\tau_h$ channel, and the observed limit of the combination are indicated as solid lines. The sensitivity of the 1-jet and 2-jet categories becomes better with increasing integrated luminosity. First, there is a large increase in the sensitivity of each jet category, which becomes weaker for larger integrated luminosities, where the systematic uncertainties dominate the overall uncertainty. Improvements on the systematic uncertainties were not taken into account because they cannot be precisely estimated. An integrated luminosity $\gtrsim 6 \text{ fb}^{-1}$ is needed at $\sqrt{s} = 13$ TeV for an upper limit which is better than the expected upper limit of the $e\tau_\mu$ channel at $\sqrt{s} = 8$ TeV. Integrated luminosities larger than 10 fb^{-1} lead to expected limits stronger than the observed upper limit of the combination at $\sqrt{s} = 8$ TeV.

8. Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$

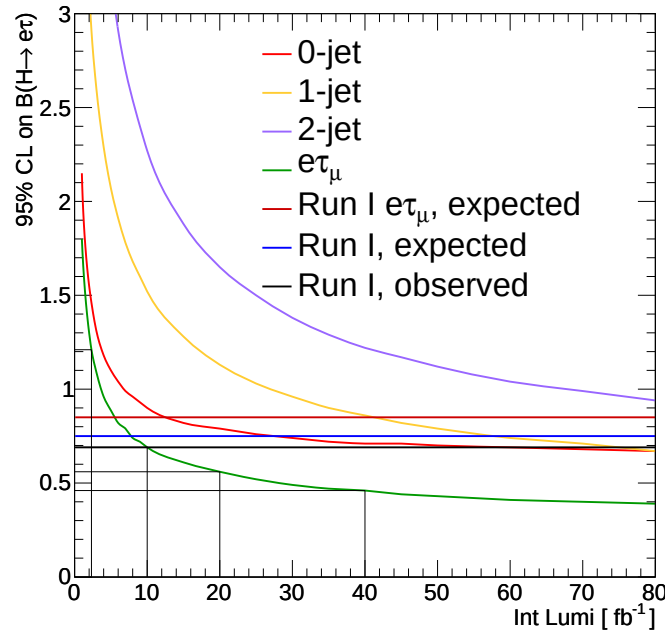


Figure 8.45.: Expected upper limit on the branching fraction $B(H \rightarrow e\tau)$ at 95% CL for $\sqrt{s} = 13$ TeV as function of the integrated luminosity. The expected limits were derived using MC simulation only. Limits are shown for each jet category, 0-jet (red), 1-jet (yellow), 2-jet (violet), and for their combination (green). Furthermore, the expected upper limits for the $e\tau_\mu$ channel (dark red) and $e\tau$ (blue) at $\sqrt{s} = 8$ TeV are indicated as line. The observed upper limit at $\sqrt{s} = 8$ TeV is indicated as black line [5].

The 2016 data set has an integrated luminosity $\sim 40 \text{ fb}^{-1}$. Data taking has ended at the end of 2016. Development of identification recipes and derivation of scale factors have not been completed in time for the presented thesis. As it can be seen in Figure 8.44, an expected upper limit of $B(H \rightarrow e\tau) \lesssim 0.5\%$ is feasible at $\sim 40 \text{ fb}^{-1}$, which would be better than the upper limit at present. This limit could be improved by an event selection which is optimised for higher integrated luminosities. New techniques to reduce the amount of $t\bar{t}$ events in the background contributions could further improve the limit.

Luminosity update In 2017, there was an update of the luminosity measurement for the 2015 collision data leading to a decrease of the integrated luminosity in comparison to the previous value. Results on the effects on the analysis are presented in Appendix I. The effects on the results of this search are very small.

In summary, the best limit on the branching fraction of $B(H \rightarrow e\tau)$ is from the search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$ at $\sqrt{s} = 8$ TeV [5], where the $e\tau_\mu$ channel has a stronger limits than the $e\tau_h$ channel, which is dominated by the misidentified-lepton background. The combined observed (expected) upper limit on $B(H \rightarrow e\tau)$ at 95% CL is 2.01% (1.32%)⁴ for the search done at $\sqrt{s} = 13$ TeV and 0.69% (0.75%) for the search done at $\sqrt{s} = 8$ TeV.

Smaller statistics than for the 8 TeV collision data lead to a reduced sensitivity due to the lower total luminosity in the 13 TeV analysis, which cannot be compensated by the increase in the cross section. A new event selection was developed to improve the expectation on the expected upper limit with respect to the event selection which was optimised for $\sqrt{s} = 8$ TeV and a total luminosity of 19.7 fb^{-1} . The cross section for $t\bar{t}$ production increases more than the other production cross sections. As this background process has the largest systematic uncertainty among the dominant background processes, a possible signal could be hidden by the large systematic uncertainties of the $t\bar{t}$ events. Thus, the reduction of the $t\bar{t}$ contributions should be more in focus of future analyses.

⁴after luminosity update, see Appendix I, previous values are: 1.94% (1.31%)

9. Summary and discussion of results

In this thesis, a search for lepton-flavour violating decays of the Higgs boson to $e\tau$ in the $e\tau_\mu$ channel was presented. A veto on jets identified as b-quark jets is applied in the search for LFV decays of the Higgs boson. The main background for identifying b-quark jets are c-quark jets because both jet types include heavy flavour hadrons. As shown in Figure 7.1, the misidentification probability depends on the type of charmed hadrons produced during hadronisation of the c-quark. Thus, a precise and accurate modelling of the production fraction is important for a precise prediction of the misidentification probability of c-quark jets as b-quark jets. Differences in the heavy flavour modelling of charmed hadrons between PYTHIA 6 and HERWIG 6 at $\sqrt{s} = 8$ TeV were studied. Production fraction, decays, and momentum fraction of charmed hadrons are differently modelled in PYTHIA 6 and HERWIG 6. These differences in heavy flavour modelling affect the probability for c-quark jets to be misidentified as b-quark jets using the Combined Secondary Vertex (CSV) identification algorithm, which combines the information of secondary vertices with track-based information.

The modelling of c-quark jets is essential in searches as the search for charged Higgs bosons which decay into an s-quark and a c-quark. Identification algorithm for c-quark jets rely on charmed hadron modelling, as they exploit the properties and signatures of charmed hadrons, such as large lifetimes, large masses, and secondary vertices. Analyses which use the identification of b-quark jets, for example to veto events with jets identified as b-quark jets, would also profit from an improved modelling of heavy-flavour hadrons which results in smaller uncertainties on the misidentification probability.

This thesis is focused on the search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$. Collision data were collected in proton-proton collisions at $\sqrt{s} = 8$ TeV and at $\sqrt{s} = 13$ TeV with the CMS experiment at the LHC. The collected data samples correspond to an integrated luminosity of 19.7 fb^{-1} and 2.3 fb^{-1} , respectively. For each data set, an optimised analysis is performed. The results of this search at $\sqrt{s} = 8$ TeV were published in Physics Letters B [5], therefore this thesis is concentrating on the search done at $\sqrt{s} = 13$ TeV.

Summary of the $e\tau_\mu$ analysis The sensitivity of the 2-jet category with respect to previous analyses [17] can be increased by taking into account Higgs production via gluon fusion in addition to Higgs production via vector boson fusion. A slightly different background method is used for the misidentified leptons because the misidentified electrons are negligible compared to misidentified muons in this analysis. However, a requirement on transverse momentum of the prompt lepton (electron) has to be applied, which is in the order of 40 GeV.

A new identification algorithm developed by the CMS collaboration is used for the muons at $\sqrt{s} = 13$ TeV, which has a lower misidentification probability. The misidentification lepton process has the highest systematic uncertainty of all background processes. Thus, possible improvements of the modelling of the misidentification background process are studied in this thesis. Control regions for different background processes are added to check the normalisation of the background processes. New discriminating variables are tested and the optimisation of the event selection is done for expected upper limit on the LFV branching fraction.

No excess of events is observed in both data samples. The best-fit branching fraction is $B(H \rightarrow e\tau) = 0.78 \left({}^{+0.22}_{-0.62} \right) \%$ for the search at $\sqrt{s} = 13$ TeV. An upper limit of $B(H \rightarrow e\tau) < 2.01\% \text{ (} 1.32\% \text{ (} {}^{+0.59}_{-0.39} \text{) expected)}$ is set at 95% CL. The corresponding search at $\sqrt{s} = 8$ TeV has an expected upper limit of $0.85 \left({}^{+0.37}_{-0.25} \right)$ in the $e\tau_\mu$ channel. A stronger observed (expected) upper combined limit of $0.69\% \text{ (} 0.75\% \text{)}$ at 95% CL was set [5]. The upper limit is dominated by the results of the $e\tau_\mu$ channel.

9.1. Comparison to other constraints

Comparison of the results to indirect limits The limits presented in this thesis are stronger as previous published indirect limits [6, 16]. Furthermore, the indirect limits are model dependent. In [6], flavour-diagonal Yukawa couplings equal to the SM values are assumed. This bound is relaxed if the flavour diagonal Yukawa couplings are smaller than their SM values. In [16], limits were calculated under the assumption of SM-like flavour-diagonal couplings, while other bounds assume the naturalness requirement that the absolute value of the mass corrections $|\delta m_l|$ is smaller than the lepton mass m_l , and a cut-off scale $\Lambda = 1$ TeV. Their bound on $\mu \rightarrow e\gamma$ implies that only $B(H \rightarrow e\tau)$ or $B(H \rightarrow \mu\tau)$ is $\mathcal{O}(10\%)$.

Search for LFV decays of the Higgs boson to $\mu\tau$ at CMS The CMS collaboration has published the first direct search for LFV decays of the Higgs boson in the $\mu\tau$ channel [17]. Collision data used in this search were collected in proton-proton collisions at $\sqrt{s} = 8$ TeV. An upper limit of $B(H \rightarrow \mu\tau) < 1.51\%$ (0.75% expected) at 95% CL was set. A small excess of collision data with a significance of 2.4σ was observed with respect to the background-only hypothesis. The best-fit branching fraction was $B(H \rightarrow \mu\tau) = 0.84^{+0.39}_{-0.37}$. Another search for LFV decays of the Higgs boson to $\mu\tau$ was done by the CMS collaboration at $\sqrt{s} = 13$ TeV [125]. No excess of events with respect to the background-only hypothesis was observed in the new data set. The $\mu\tau_e$ channel has an observed upper limit of $B(H \rightarrow \mu\tau) < 1.15\%$ (1.96% expected). An upper limit of $B(H \rightarrow \mu\tau) < 1.20\%$ (1.26% expected) was set at 95% CL for the combination with the $\mu\tau_h$ channel.

Comparison of the $e\tau_\mu$ analysis to the $\mu\tau_e$ analysis As the $e\tau_\mu$ and $\mu\tau_e$ channels have a similar signature, only the prompt lepton and the lepton of the tau decay are swapped, the sensitivity is expected to be similar, too. At $\sqrt{s} = 8$ TeV, the sensitivity is comparable comparing the expected limits of each jet category. The $\mu\tau_e$ channel has a better sensitivity for the 0-jet category, while the $e\tau_\mu$ has better sensitivities in the 1-jet and 2-jet categories. A weaker vector-boson fusion selection is applied in the 2-jet category of the $e\tau_\mu$ analysis to enhance the signal.

At $\sqrt{s} = 13$ TeV, the $e\tau_\mu$ analysis has a better sensitivity than the $\mu\tau_e$ analysis as the event selection was optimised for the reduced statistics. In addition, the sensitivity of the $e\tau_\mu$ analysis is comparable to the sensitivity of the combined analysis $\mu\tau$ at $\sqrt{s} = 13$ TeV.

The sensitivity of the searches for LFV decays of the Higgs boson are better using the data sample of 19.7 fb^{-1} collected in proton-proton collisions at $\sqrt{s} = 8$ TeV than the data sample of 2.3 fb^{-1} collected at $\sqrt{s} = 13$ TeV. The reason for this is the reduced statistics, as the increase in the cross sections cannot compensate the smaller integrated luminosity. Optimising the event selection to the reduced statistics, results in a better sensitivity.

Search for LFV decays of the Higgs boson to $e\mu$ at CMS The CMS collaboration has also published a direct search for LFV decays of the Higgs boson to $e\mu$ using a data sample collected in proton-proton collisions at $\sqrt{s} = 8$ TeV [5]. An observed limit of $B(H \rightarrow e\mu) < 0.035\%$ (< 0.048 expected) was set. This is weaker than the indirect limit $B(H \rightarrow e\mu) < 3 \times 10^{-9}$ [16], which holds under the hypothesis of SM-like flavour-diagonal couplings of the Higgs boson.

Search for LFV decays of the Higgs boson at ATLAS The ATLAS collaboration has published direct searches for lepton-flavour violating decays of the Higgs boson to $e\tau$ and $\mu\tau$ [147, 148]. A data sample with an integrated luminosity of 20.3 fb^{-1} was collected in proton-proton collisions at $\sqrt{s} = 8 \text{ TeV}$. Each search was performed in a leptonic and a hadronic channel. The searches in the hadronic channels are model independent, while leptonic channels [148] are based on the assumption that either the Yukawa couplings for the decay of the Higgs boson to $e\tau$ or the one for the decay to $\mu\tau$ is zero. Both channels are discussed in Appendix J.

No significant excess was found and observed upper limits of $B(H \rightarrow \mu\tau) < 1.85$ ($1.24^{+0.50}_{-0.35}$ expected) and $B(H \rightarrow e\tau) < 1.81$ ($2.07^{+0.82}_{-0.58}$ expected) were set using the hadronic channels. Model dependent observed upper limits of $B(H \rightarrow \mu\tau) < 1.79$ ($1.73^{+0.74}_{-0.49}$ expected) and $B(H \rightarrow e\tau) < 1.36$ ($1.48^{+0.60}_{-0.42}$ expected) are derived using the leptonic channels. The model-independent hadronic channels and the model-dependent leptonic channels are combined [148]. Model dependent observed upper limits of $B(H \rightarrow \mu\tau) < 1.43$ ($1.01^{+0.40}_{-0.29}$ expected) and $B(H \rightarrow e\tau) < 1.04$ ($1.21^{+0.49}_{-0.34}$ expected) are derived.

The sensitivity of the ATLAS combination is weaker than the corresponding combinations of the CMS analyses at $\sqrt{s} = 8 \text{ TeV}$. Observed limits set with the CMS analyses are better than the ATLAS ones, when including the CMS $\mu\tau$ analysis at $\sqrt{s} = 13 \text{ TeV}$. In addition, only the limits from the hadronic channels are model independent.

Summary of upper limits on LFV decays of the Higgs boson A summary of the observed and expected upper limits on the lepton-flavour violating branching fractions of the Higgs boson is shown in Figure 9.1. Only model independent LHC results are shown. The SM branching fractions of the Higgs boson are indicated as lines. Assuming a similar behaviour proportional to the involved lepton masses, the branching fractions of LFV Higgs boson decays are shown for the naturalness assumption, discussed in Section 3.1. The LHC limits on the LFV decays of the Higgs boson involving taus are close to the naturalness limit of the Higgs boson decay to $\mu\tau$, while the LHC limit on $B(H \rightarrow e\mu)$ is close to the SM branching fraction of $H \rightarrow \mu\mu$. Only the limits on $B(H \rightarrow \mu\tau)$ are close to the corresponding naturalness assumption, while limits on the other two branching fractions are about two orders of magnitude weaker than these assumptions. Thus, a much larger sensitivity of the corresponding analyses is required to set limits or to measure the branching fractions close to their natural branching fraction.

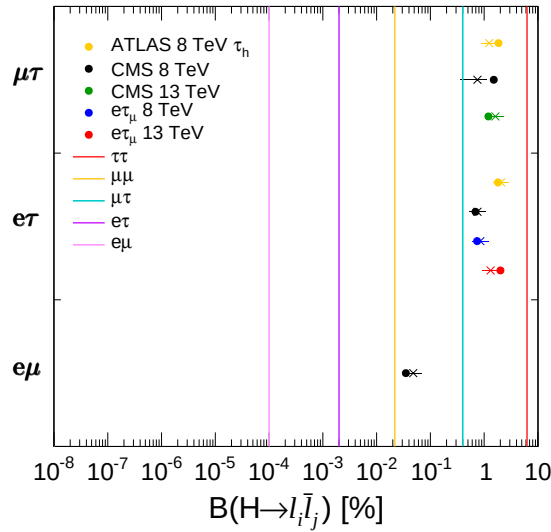
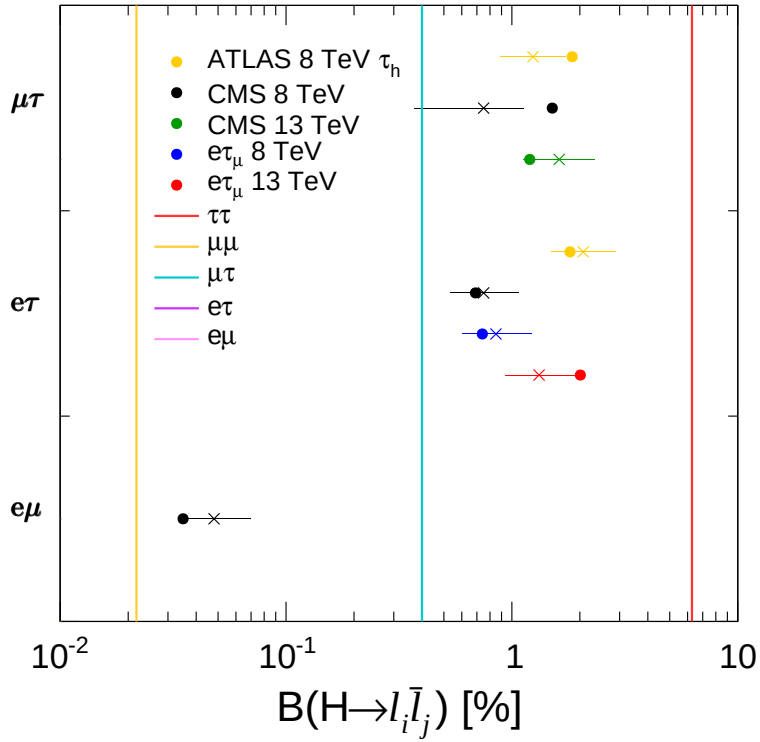


Figure 9.1.: Summary of the upper limits on the branching fraction of LFV Higgs boson decays. The observed (circle) and expected (cross) upper limits at 95% on the branching fractions of the LFV Higgs-boson decays are shown. The one standard deviations of the expected upper limit are shown as solid line. Data samples were collected in proton-proton collisions at $\sqrt{s} = 8$ TeV with the CMS experiment (CMS 8 TeV, black) [5, 17] and at the ATLAS experiment (ATLAS 8 TeV τ_h , yellow). Another data sample was collected in proton-proton collisions at $\sqrt{s} = 13$ TeV with the CMS experiment (CMS 13 TeV, green) [125]. Results of this thesis are indicated as $e\tau_\mu$ 8 TeV (blue) and as $e\tau_\mu$ 13 TeV (red). The SM Higgs branching fractions to leptons [26, 27] as well as the naturalness assumption on the LFV branching fractions are indicated as solid lines, where it is assumed that the LFV branching fractions are proportional to the involved lepton masses. A summary of the observed and expected LHC limits and expected SM values is shown in the bottom figure. A zoom to the LHC limits is shown in the top figure.

9. Summary and discussion of results

The best upper limits on the branching fractions of LFV decays of the Higgs boson come from the CMS experiment, when considering only direct limits. For LFV decays of the Higgs boson to a muon and a tau, the best observed (expected) upper limit is $B(H \rightarrow \mu\tau) < 1.20\%$ (1.26% expected) at 95% CL from the analysis at $\sqrt{s} = 13$ TeV [125]. For LFV decays of the Higgs boson to an electron and a tau, the best observed (expected) upper limit is 0.69% (0.75%) at 95% CL from the analysis at $\sqrt{s} = 8$ TeV [5]. The direct limit from the CMS experiment is weaker than the indirect limits [6]. However, they are model independent. An observed limit of $B(H \rightarrow e\mu) < 0.035\%$ (< 0.048 expected) at 95% CL could be set at $\sqrt{s} = 8$ TeV.

9.2. Possible improvements

As possible improvements of the presented analyses, kinematic fits on the Higgs and Z boson, and the top-quark mass could be done. The χ^2 of the fits can be used to reduce the SM background processes. Multivariate analysis techniques could be used to optimise the selection of the events taking correlations of the observables into account. Boosted decision trees (BDTs) can be trained for different SM background processes. Their results could be used to classify the events into several categories, i.e. to classify an event as signal-like, as top background or as Drell-Yan background event. Then, only events classified as signal could be selected and the normalisation for each background process could be checked individually using the other categories.

Appendices

A. Mass resolution

A.1. Derivation of the collinear mass M_{col} reconstruction for $H \rightarrow e\tau_\mu$

The Higgs boson (H) decays into an electron (e) and a tau (τ). Thus, the mass of the Higgs boson depends on the energy of the electron, the tau energy, and the angle between both leptons $\alpha_{e,\tau}$.

$$\begin{aligned} M_H^2 &= (p_e + p_\tau)^2 \\ &= m_e^2 + m_\tau^2 + 2(E_e E_\tau - \vec{p}_e \vec{p}_\tau) \\ &\approx 2E_e E_\tau (1 - \cos \alpha_{e,\tau}), \end{aligned} \tag{A.1}$$

when $m_e \ll |\vec{p}_e|$ and $m_\tau \ll |\vec{p}_\tau|$.

The tau decays further into a muon (μ) and corresponding neutrinos (ν) which lead to missing energy. Then, the visible fraction of the tau decay is defined by the ratio of the muon momentum to the momentum of the tau.

$$x_\tau^{\text{vis}} \equiv \frac{|\vec{p}_\mu|}{|\vec{p}_\tau|} = \frac{|\vec{p}_\mu|}{|\vec{p}_\mu| + |\vec{p}_\nu|} \tag{A.2}$$

Using Equation A.2, the equation for the energy of the tau as function of mass and momentum can be transformed as function of muon energy and visible fraction of the tau decay.

$$\begin{aligned} E_\tau &= \sqrt{m_\tau^2 + |\vec{p}_\tau|^2} \approx \sqrt{|\vec{p}_\tau|^2} \\ &= |\vec{p}_\tau| = |\vec{p}_\mu| + |\vec{p}_\nu| \\ &= \frac{|\vec{p}_\mu| + |\vec{p}_\nu|}{|\vec{p}_\mu|} \cdot |\vec{p}_\mu| \\ &= \frac{1}{x_\tau^{\text{vis}}} \cdot |\vec{p}_\mu| \\ &\approx \frac{1}{x_\tau^{\text{vis}}} \cdot E_\mu \end{aligned} \tag{A.3}$$

A.1. Derivation of the collinear mass M_{col} reconstruction for $H \rightarrow e\tau_\mu$

In Equation A.1, the energy of the tau can be replaced as ratio of the muon energy to the visible energy fraction. Furthermore, the collinear approximation is used, which states that the tau decay-products are collinear to the tau direction, therefore the neutrinos are collinear to the muon direction.

$$\begin{aligned} 2E_e \frac{E_\mu}{x_\tau^{\text{vis}}} (1 - \cos \alpha_{e,\mu}) &\approx \frac{1}{x_\tau^{\text{vis}}} \cdot M_{\text{vis}}^2 \\ &\equiv M_{\text{col}}^2 \end{aligned} \quad (\text{A.4})$$

The term with the energy of the electron and muon energy describes the visible invariant-mass. Thus, the collinear mass is given by the visible mass and the square root of the visible fraction from the tau decay.

$$M_{\text{col}} = \frac{M_{\text{vis}}}{\sqrt{x_\tau^{\text{vis}}}} \quad (\text{A.5})$$

As only the transverse energy of a collision is known, this quantity has to be calculated in the transverse plane. Neutrinos are reconstructed as missing transverse energy. Using only the part of $\not{E}_T$ which is collinear to the muon direction, the projected missing transverse energy $\not{E}_{T,\text{proj}}$ approximates the energy from the neutrinos in the transverse plane.

$$p_T^\nu \approx \not{E}_{T,\text{proj}} = \frac{\vec{\not{E}}_T \cdot \vec{p}_T^{\tau^{\text{vis}}}}{|\vec{p}_T^{\tau^{\text{vis}}}|} \quad (\text{A.6})$$

Furthermore, the visible fraction of the tau decay can be expressed in terms of the visible momentum of the tau decay and $\not{E}_{T,\text{proj}}$.

$$x_\tau^{\text{vis}} = \frac{p_T^{\tau^{\text{vis}}}}{p_T^{\tau^{\text{vis}}} + \not{E}_{T,\text{proj}}} \quad (\text{A.7})$$

A. Mass resolution

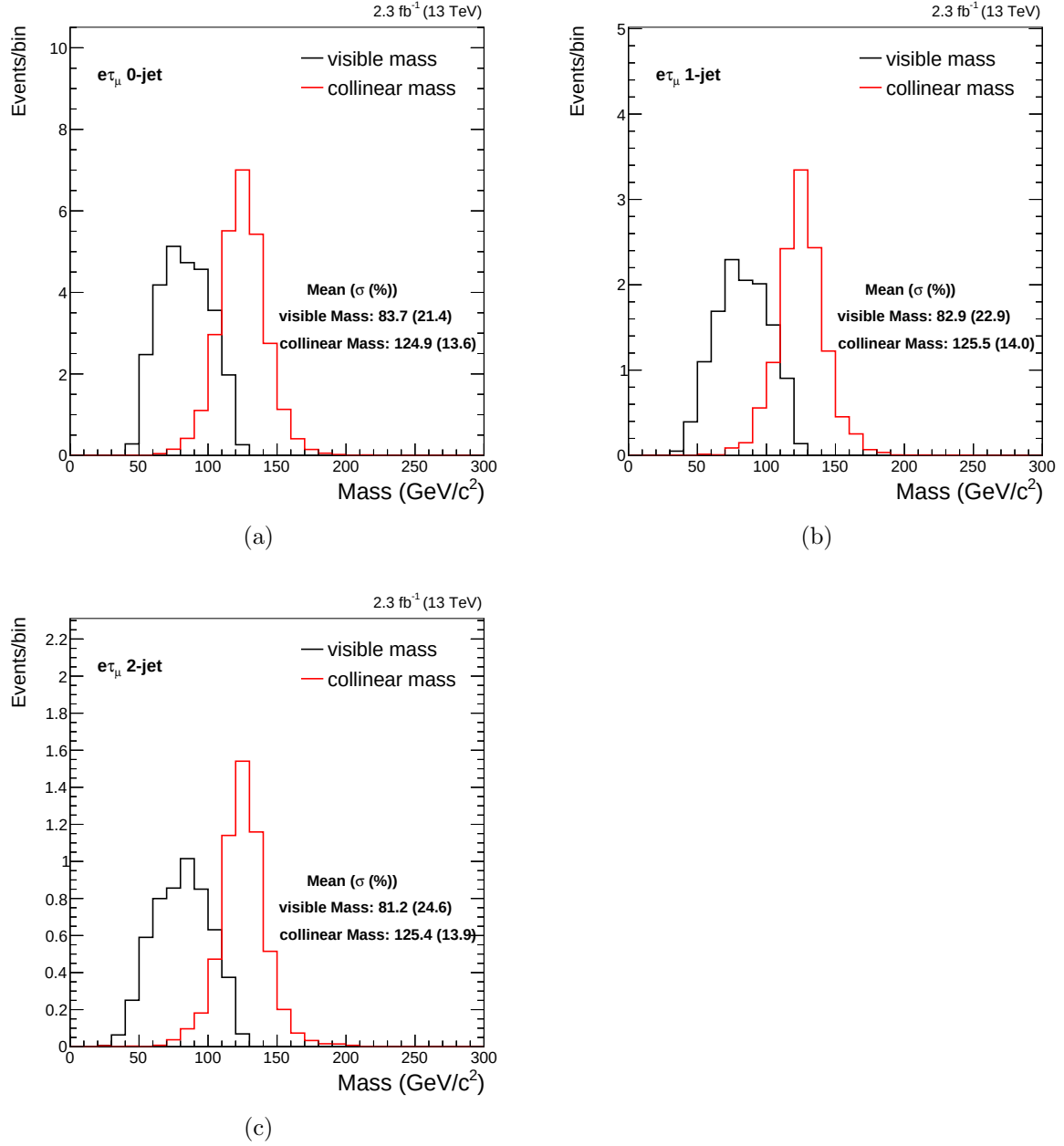


Figure A.1.: Visible mass (black line), and collinear mass (red line) are shown for LFV decays of the Higgs boson to $e\tau\mu$ for events with (a) no jet, (b) one jet, or (c) two jets with $p_T > 30$ GeV.

A.1. Derivation of the collinear mass M_{col} reconstruction for $H \rightarrow e\tau_\mu$

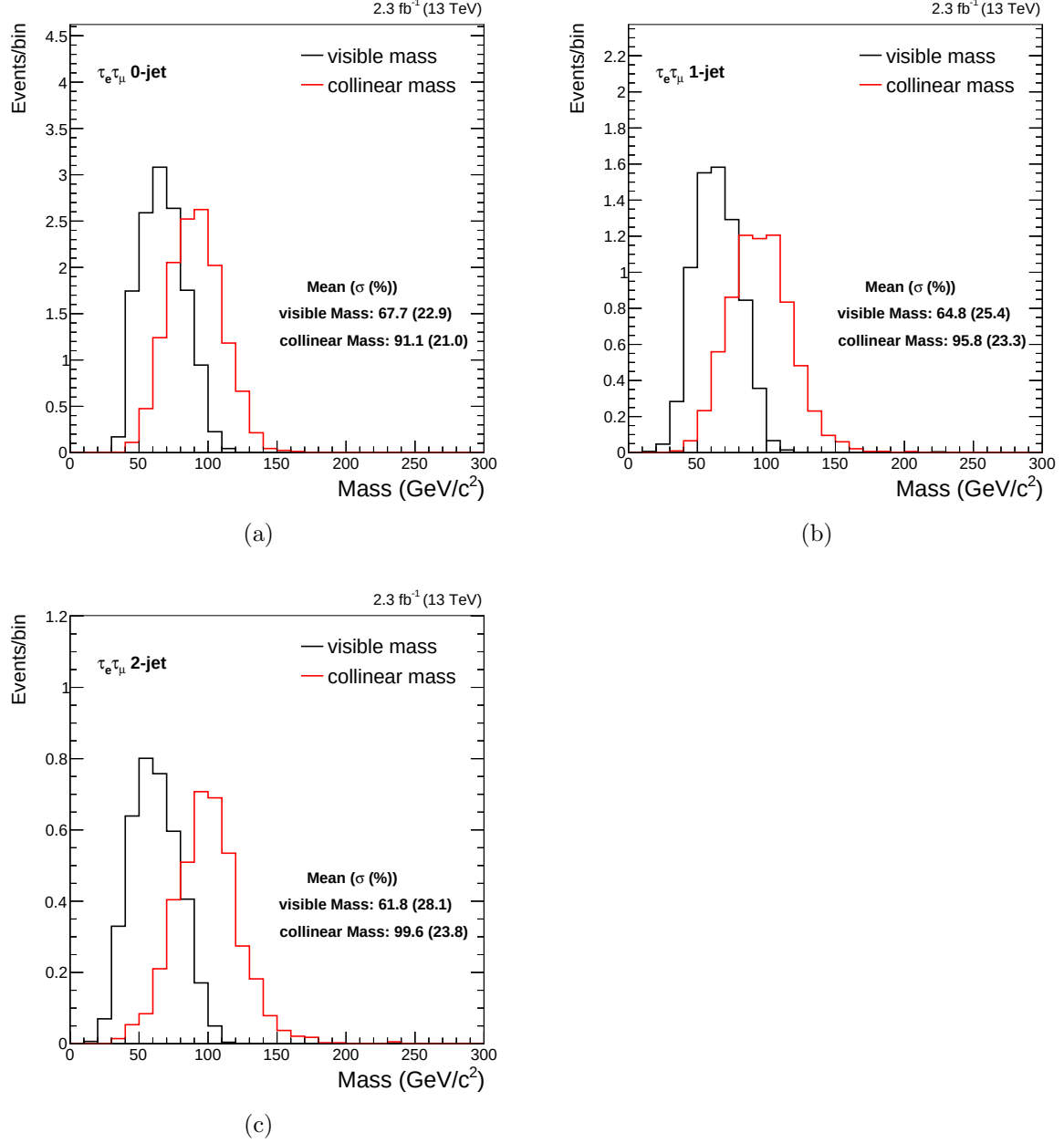


Figure A.2.: Visible mass (black line), and collinear mass (red line) are shown for SM decays of the Higgs boson to $\tau_e\tau_\mu$ for events with (a) no jet, (b) one jet, or (c) two jets with $p_T > 30$ GeV.

B. Event simulation

B. Event simulation

Table B.1.: Overview of MC samples used in the LFV $H \rightarrow e\tau_\mu$ analysis. Events were generated for proton-proton collisions with a centre-of-mass energy of $\sqrt{s} = 13$ TeV and a 25 ns bunch spacing using PYTHIA 8 and CUETP8M1 tune. The detector response was simulated using GEANT4.

Process	Remark	Matrix element	MadSpin	Matching/ Merging	Matrix element/ Hadronisation	Decays
$H \rightarrow e\tau_\mu$ (GF)	$M_H = 125$ GeV	POWHEG	-	-	PYTHIA 8	-
$H \rightarrow e\tau_\mu$ (VBF)	$M_H = 125$ GeV	POWHEG	-	-	PYTHIA 8	-
$H \rightarrow \tau\tau$ (GF)	$M_H = 125$ GeV	POWHEG	-	-	PYTHIA 8	-
$H \rightarrow \tau\tau$ (VBF)	$M_H = 125$ GeV	POWHEG	-	-	PYTHIA 8	-
$DY \rightarrow ll$	$M_{ll} > 50$ GeV	MadGraph	-	MLM	PYTHIA 8	-
$DY \rightarrow ll$	$M_{ll} > 50$ GeV	aMC@NLO	-	FxFx	PYTHIA 8	-
$W + \text{jets} \rightarrow l\nu$	inclusive	MadGraph	-	MLM	PYTHIA 8	-
$t\bar{t}$	inclusive	POWHEG	-	-	PYTHIA 8	-
$t\bar{t}$	inclusive	POWHEG	-	-	PYTHIA 8	EvtGen
$tW/\bar{t}W$	inclusive	POWHEG	-	-	PYTHIA 8	-
WW	inclusive	-	-	-	PYTHIA 8	-
WZ	inclusive	-	-	-	PYTHIA 8	-
ZZ	inclusive	-	-	-	PYTHIA 8	-
$WW \rightarrow 2l2\nu$	-	POWHEG	-	-	PYTHIA 8	-
$WZ \rightarrow 2l2q$	-	aMC@NLO	✓	FxFx	PYTHIA 8	-
$ZZ \rightarrow 2l2q$	-	aMC@NLO	✓	FxFx	PYTHIA 8	-
$ZZ \rightarrow 2l2\nu$	-	POWHEG	-	-	PYTHIA 8	-
$W\gamma \rightarrow l\nu\gamma$	-	MadGraph	-	MLM	PYTHIA 8	-
$W\gamma^* \rightarrow l\nu ee$	0,1,2 jets	MadGraph	-	MLM	PYTHIA 8	-
$W\gamma^* \rightarrow l\nu\mu\mu$	0,1,2 jets	MadGraph	-	MLM	PYTHIA 8	-

Table B.2.: Cross sections, and cross-section accuracy for simulated samples used in the LFV analysis.

Process	cross section [pb]	Accuracy
$H \rightarrow e\tau$ (GF)	$44.14^{+7.6}_{-8.1}$ (QCD scale) ± 3.1 (PDF+ α_S) pb	NNLO+NNLL QCD and NLO EW [142]
$H \rightarrow e\tau$ (VBF)	$3.782^{+0.4}_{-0.3}$ (QCD scale) ± 2.1 (PDF+ α_S) pb	NNLO QCD and NLO EW [142]
$H \rightarrow \tau\tau$ (GF)	$44.14^{+7.6}_{-8.1}$ (QCD scale) ± 3.1 (PDF+ α_S) pb	NNLO+NNLL QCD and NLO EW [142]
$H \rightarrow \tau\tau$ (VBF)	$3.782^{+0.4}_{-0.3}$ (QCD scale) ± 2.1 (PDF+ α_S) pb	NNLO QCD and NLO EW [142]
$DY \rightarrow ll$	5765.4 ± 1.8 (integration) ± 99.6 (PDF+ α_S)	NNLO
$W+\text{jets} \rightarrow l\nu$	$61526.7^{+497.1}_{-264.6} \pm 2312.7$	NNLO
$t\bar{t}$	772	measured [129]
$tW/\bar{t}W$	$35.85 \pm 0.90 \pm 1.70$	approx. NNLO
WW	$118.7^{+2.5\%}_{-2.2\%}$ pb	NNLO [130]
WZ	47.13	NLO
ZZ	16.523	NLO
$WW \rightarrow 2l2\nu$	12.178	NNLO
$WZ \rightarrow 2l2q$	5.595	NLO
$ZZ \rightarrow 2l2q$	3.22	NLO
$ZZ \rightarrow 2l2\nu$	0.564	NLO
$W\gamma \rightarrow l\nu\gamma$	405.271	LO
$W\gamma^* \rightarrow l\nu ee$	3.526	LO
$W\gamma^* \rightarrow l\nu\mu\mu$	2.793	LO

C. Misidentified lepton background

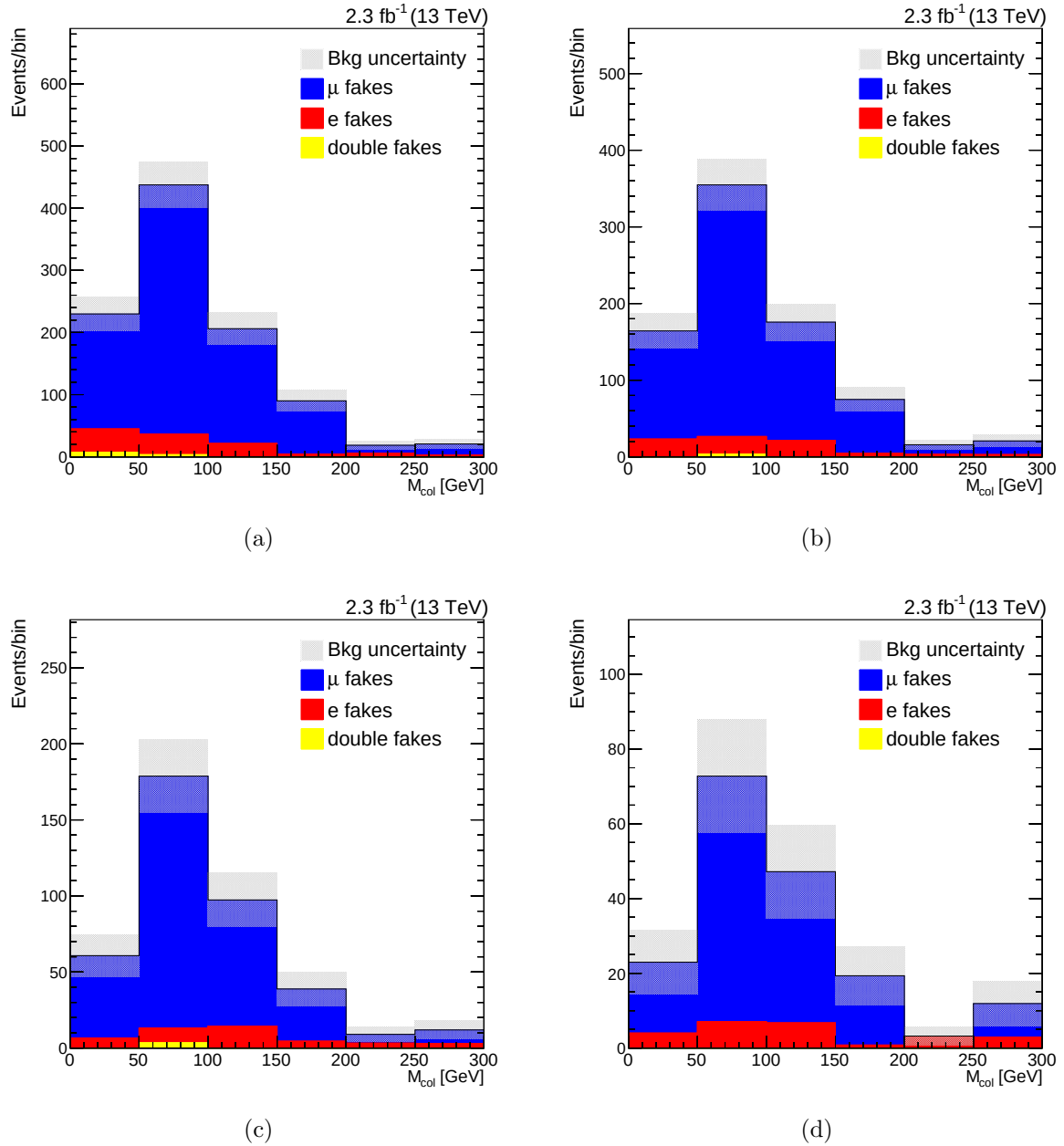


Figure C.1.: Composition of the misidentified lepton background after preselection as function of the collinear mass for (a) $p_T^e > 20$ GeV, (b) $p_T^e > 30$ GeV, (c) $p_T^e > 40$ GeV, and (d) $p_T^e > 50$ GeV.

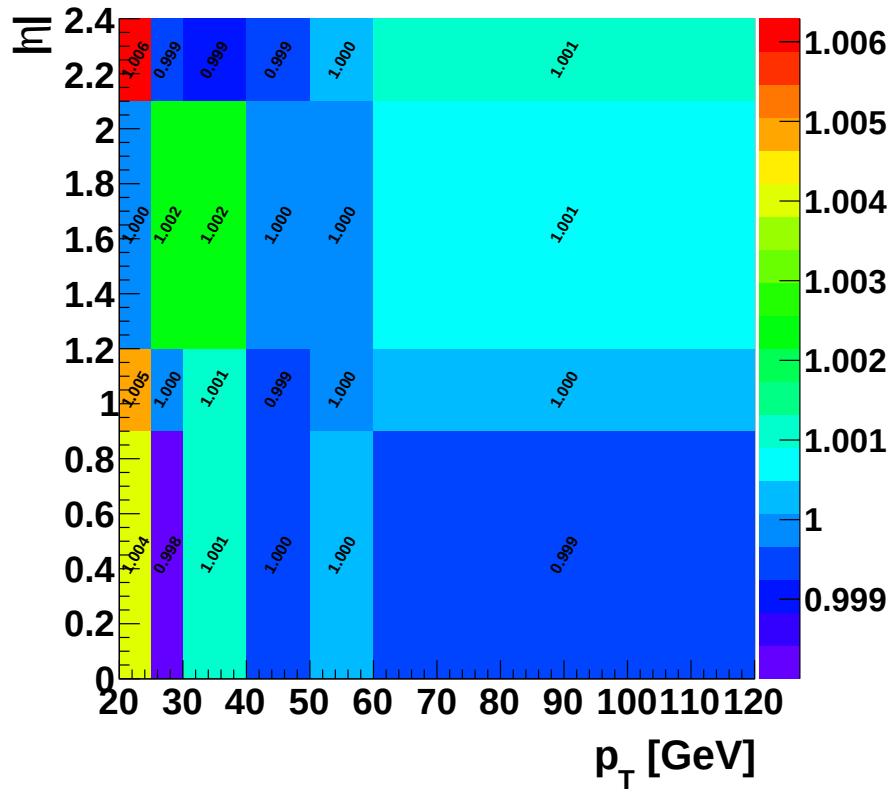


Figure C.2.: Muon scale-factors for the loose-isolation requirement as function of muon p_T and $|\eta|$. These scale factors were centrally computed by the CMS collaboration using a tag-and-probe method and the medium muon ID.

C. Misidentified lepton background

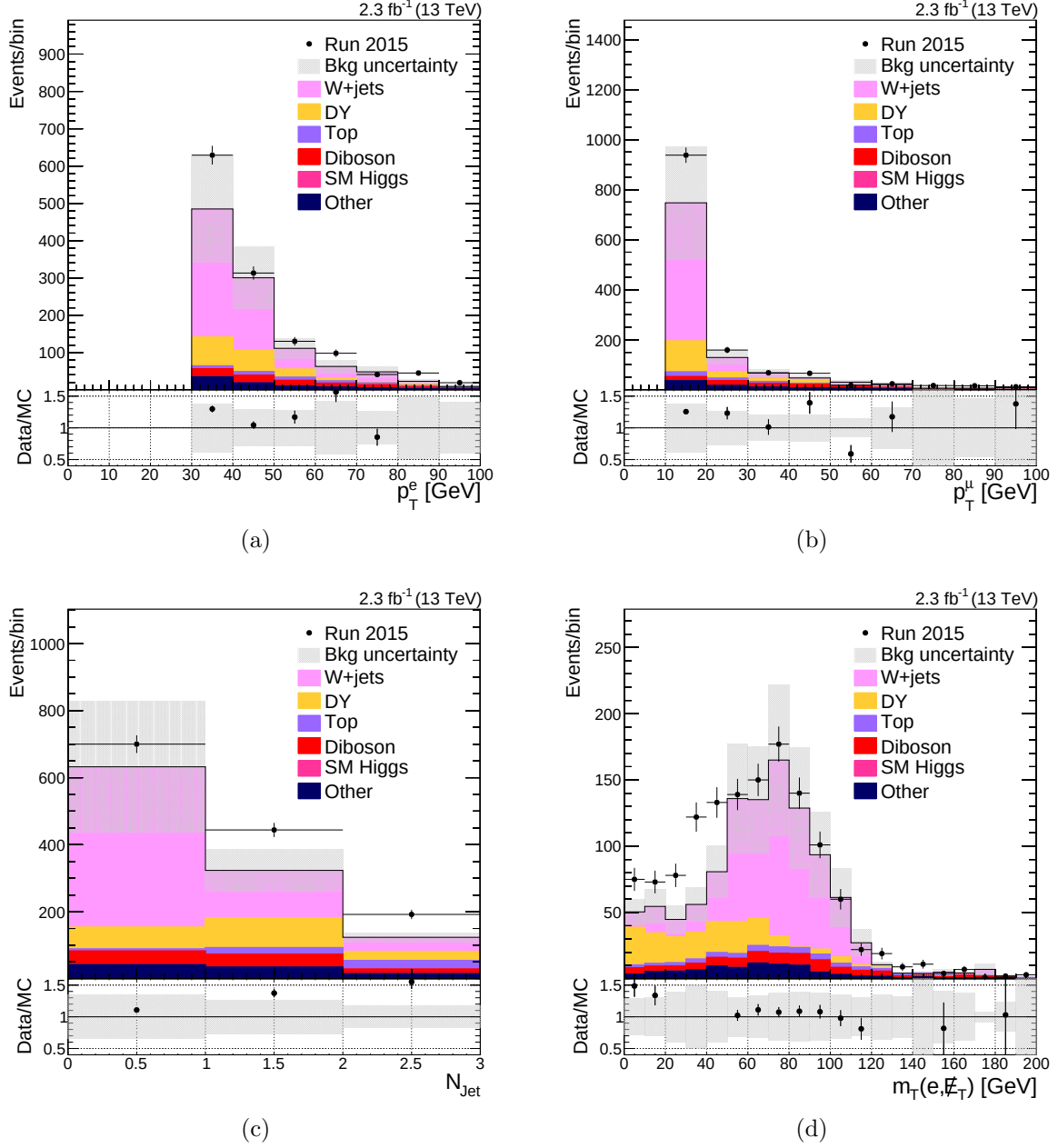


Figure C.3.: Validation of the misidentified-lepton background estimated using MC simulation. Distributions are shown after preselection in the same-charge region: (a) p_T^e , (b) p_T^μ , (c) N_{Jet} , and (d) $m_T(e, \cancel{E}_T)$.

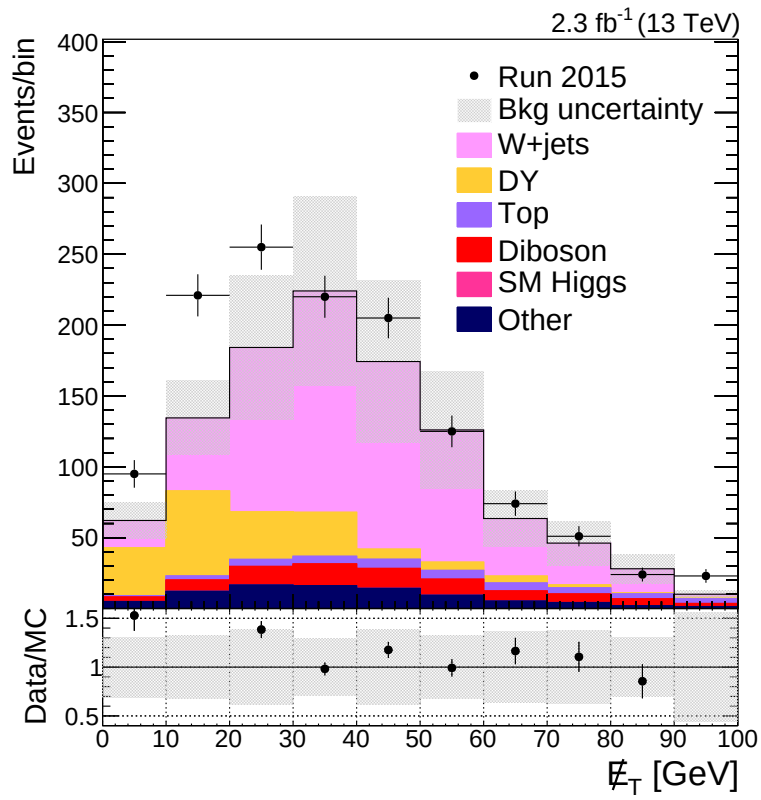


Figure C.4.: Event yield of the misidentified lepton-background in the same-charge region as function of E_T . Events were generated using PYTHIA 8.

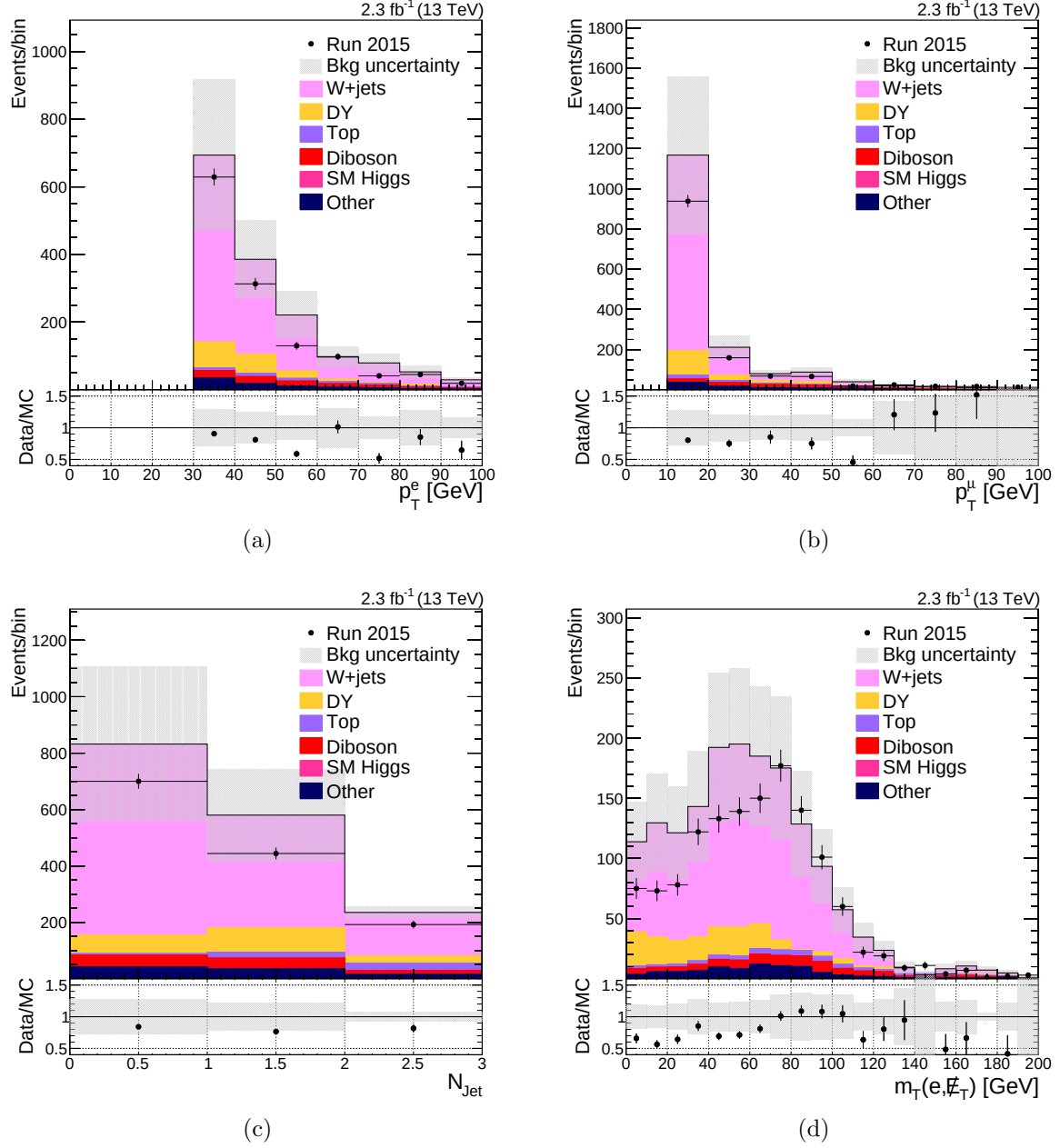


Figure C.5.: Validation distributions of the isolation side-band method are shown: (a) p_T^e , (b) p_T^μ , (c) N_{Jet} , and (d) $m_T(e, \cancel{E}_T)$. These distributions are shown after preselection for the same-charge region II using the isolation side-band method for the estimation of the misidentified-lepton background.

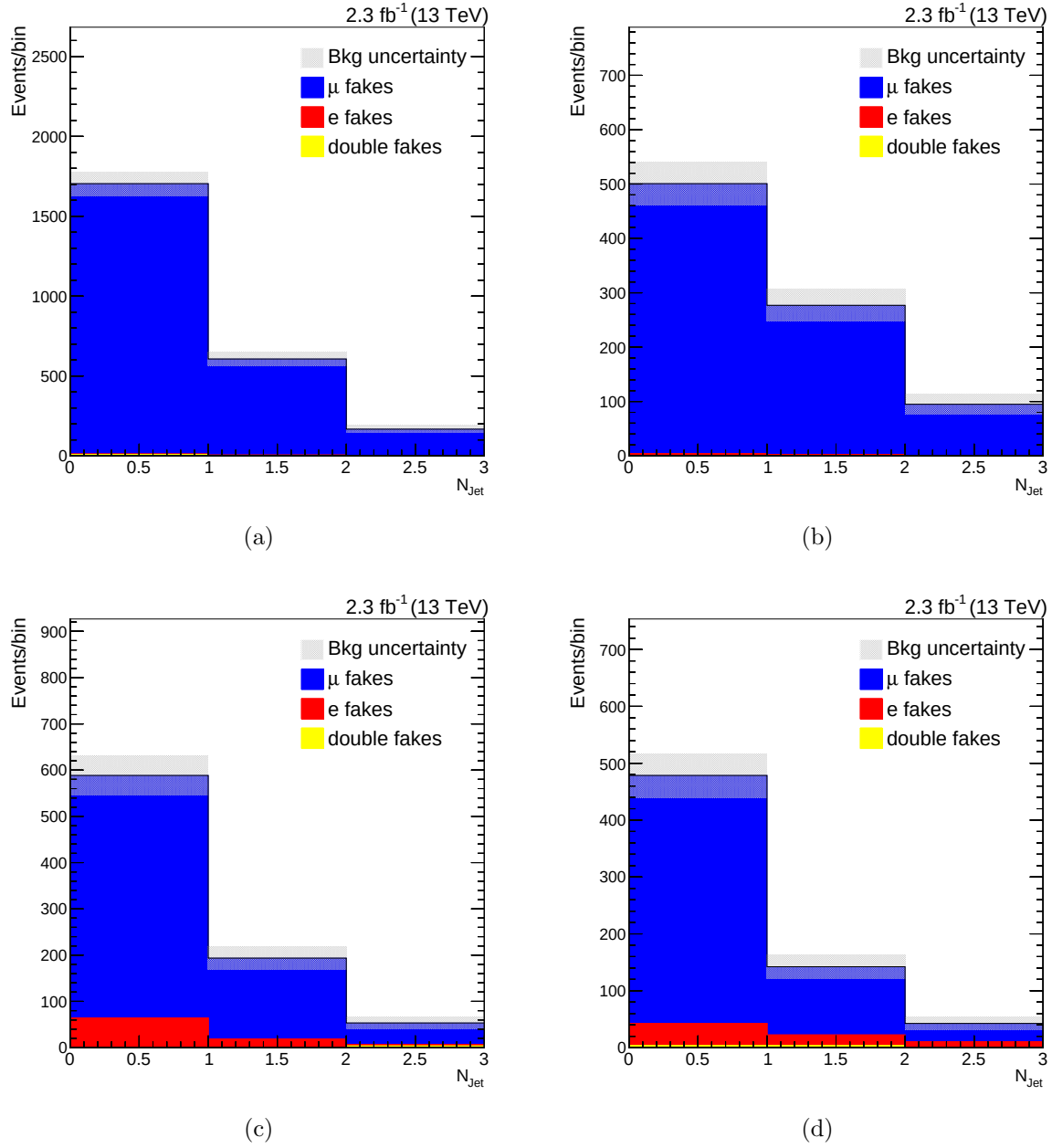


Figure C.6.: Composition of the misidentified-lepton background as function of the jet multiplicity after preselection in (a) region III, (b) region IV, (c) region I, and (d) region II. Events were generated using MC simulation.

C. Misidentified lepton background

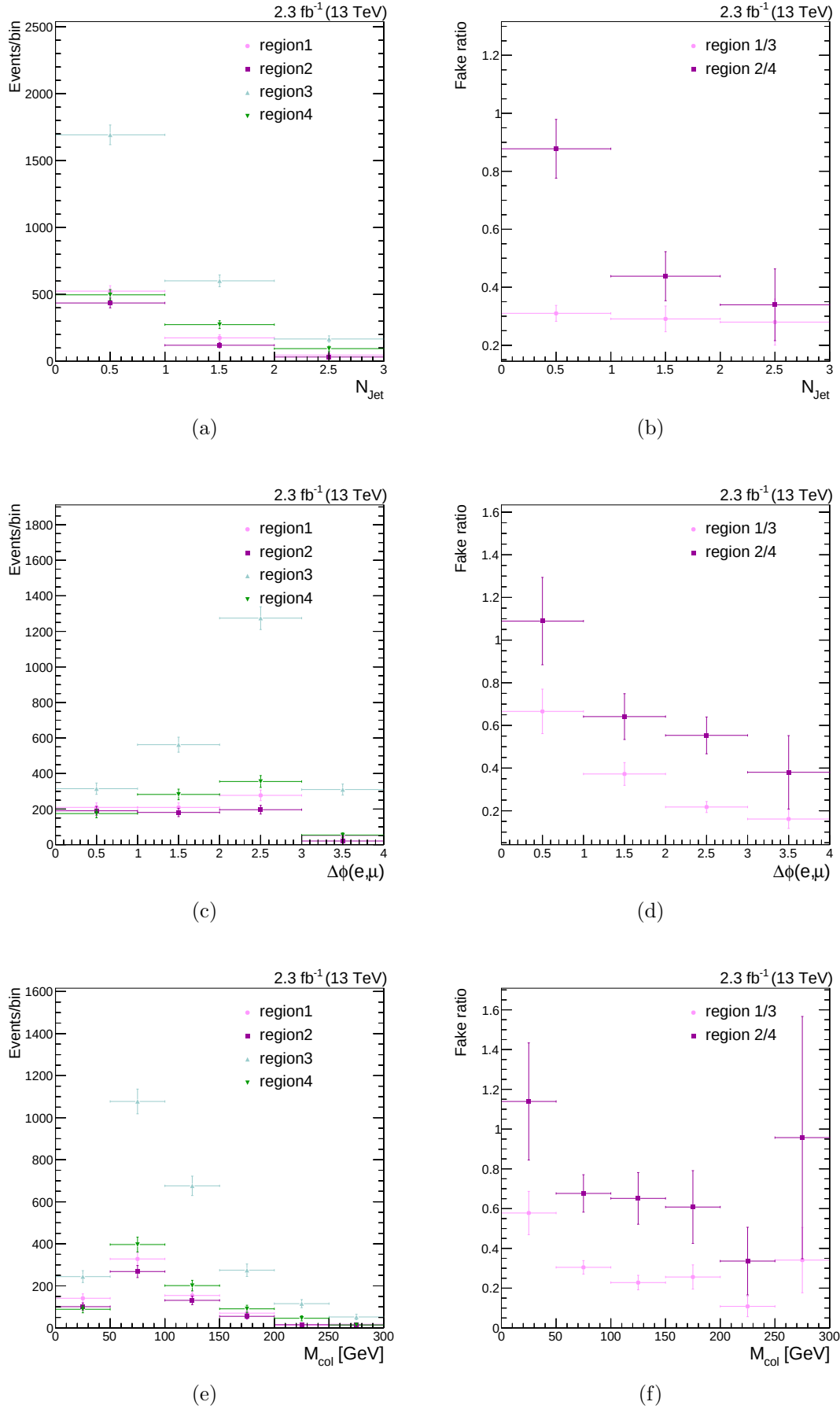


Figure C.7.: The misidentified-muon background is shown in each region for simulation after preselection. Number of events are shown as function of (a) N_{Jet} , (c) $\Delta\phi(e, \mu)$, and (e) M_{col} . Furthermore, the misidentification ratio is shown for opposite-charge and same-charge leptons as function of (b) N_{Jet} , (d) $\Delta\phi(e, \mu)$, and (f) M_{col} .

D. Control regions

D. Control regions

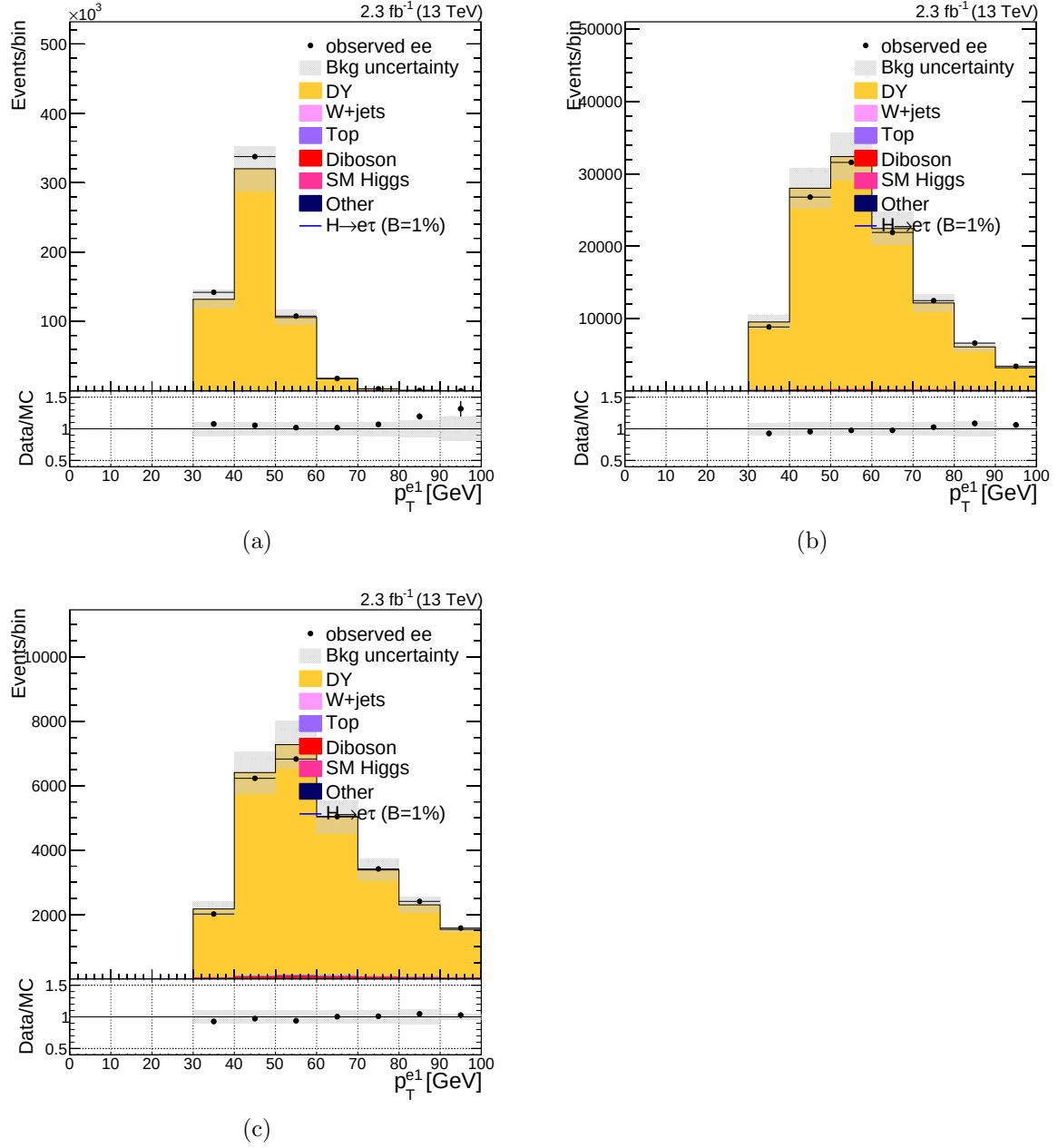


Figure D.1.: Drell-Yan background in the Drell-Yan control region as function of electron p_T for the electron which has the largest p_T . The background process is shown for each jet category: (a) 0-jet category, (b) 1-jet category, and (c) 2-jet category. Collision events were collected using a single electron trigger and the 2015 data set. Events were generated with aMC@NLO applying a systematic uncertainty of 10%.

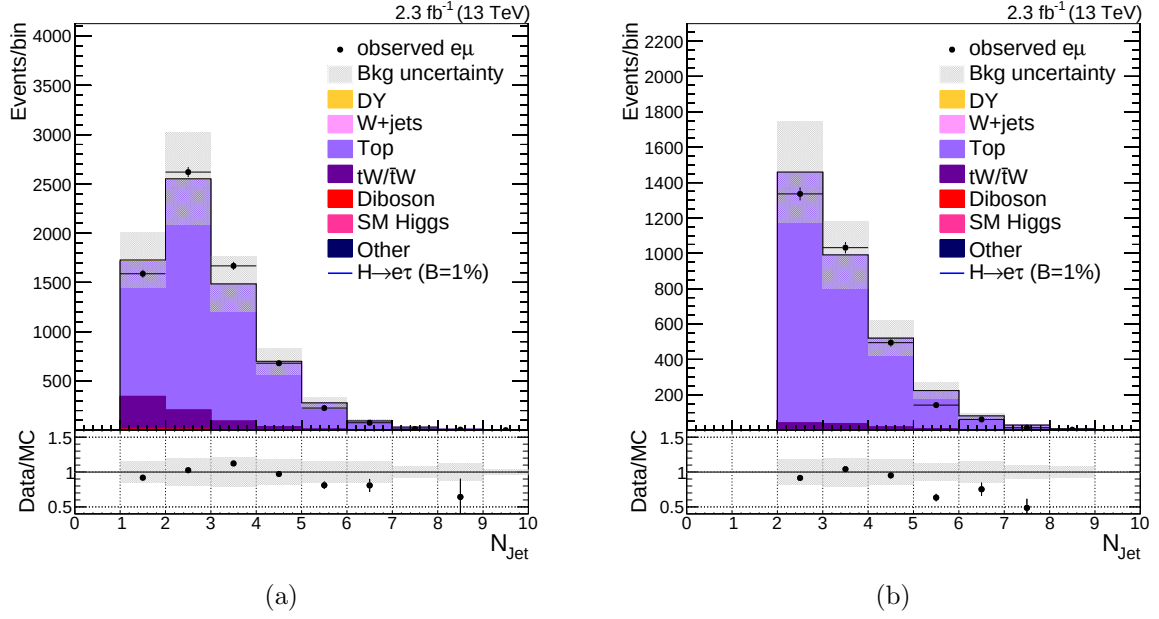


Figure D.2.: Validation of the $t\bar{t}$ process in the top-background control-region. The event yield is shown as function of the jet multiplicity for (a) 1 identified b-quark jet, and (b) 2 identified b-quark jets. Events were generated with POWHEG for the matrix-element generation and PYTHIA 8 for the parton shower. Simulated $t\bar{t}$ events were normalised according to the measured production cross-section [129].

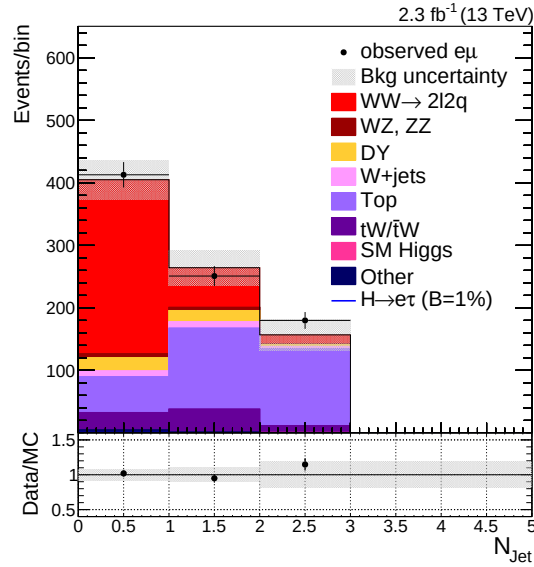


Figure D.3.: Event yield as function of the jet multiplicity in the W^+W^- process control-region. Events were generated with POWHEG for the matrix-element generation and PYTHIA 8 for the parton shower. A calculated inclusive cross section was used for W^+W^- production [130].

D. Control regions

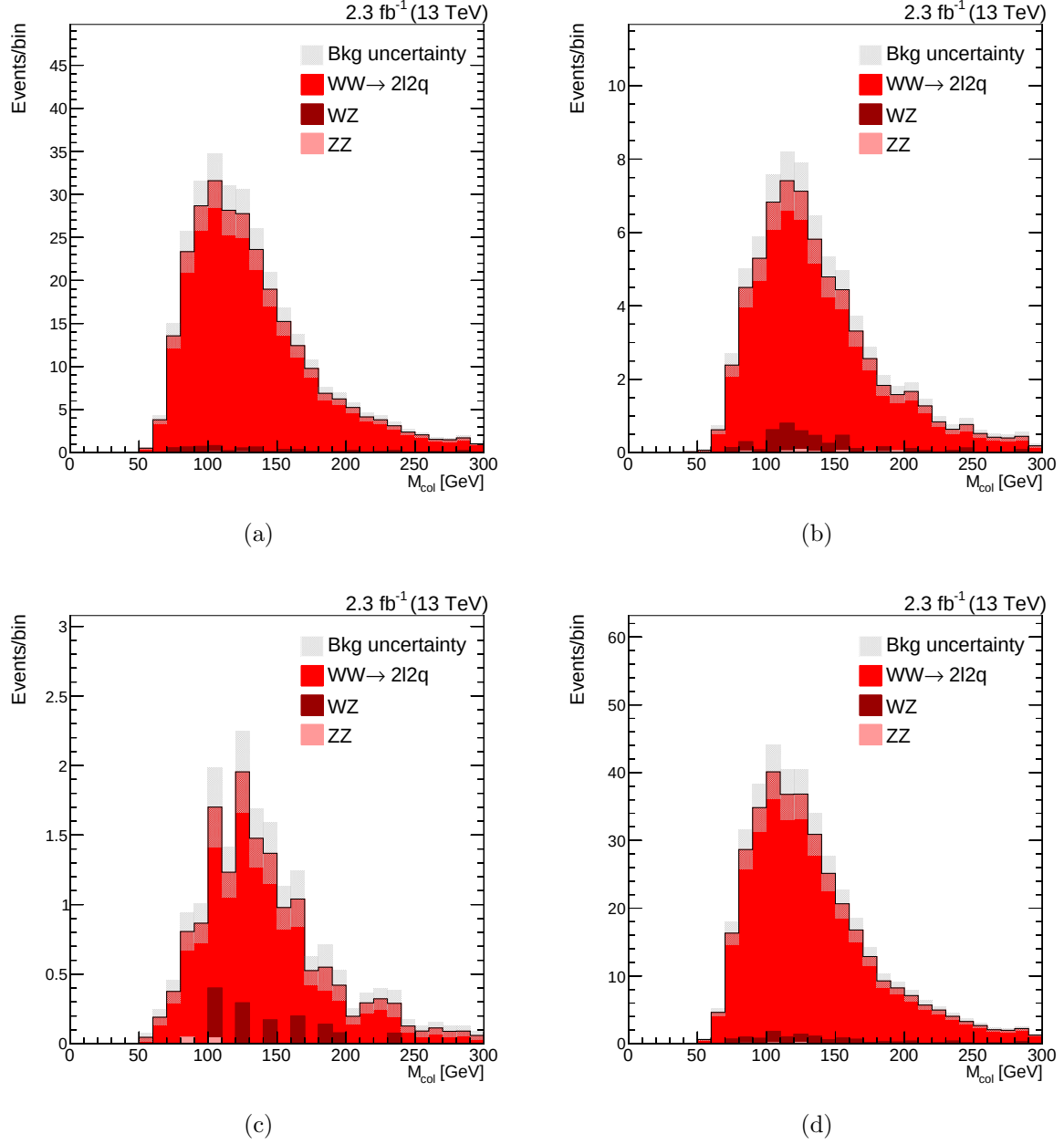


Figure D.4.: Composition of the diboson background after preselection as function of collinear mass for the (a) 0-jet, (b) 1-jet, (c) 2-jet category, and (d) for all three jet categories combined. A uncertainty of 10% was used for each diboson process.

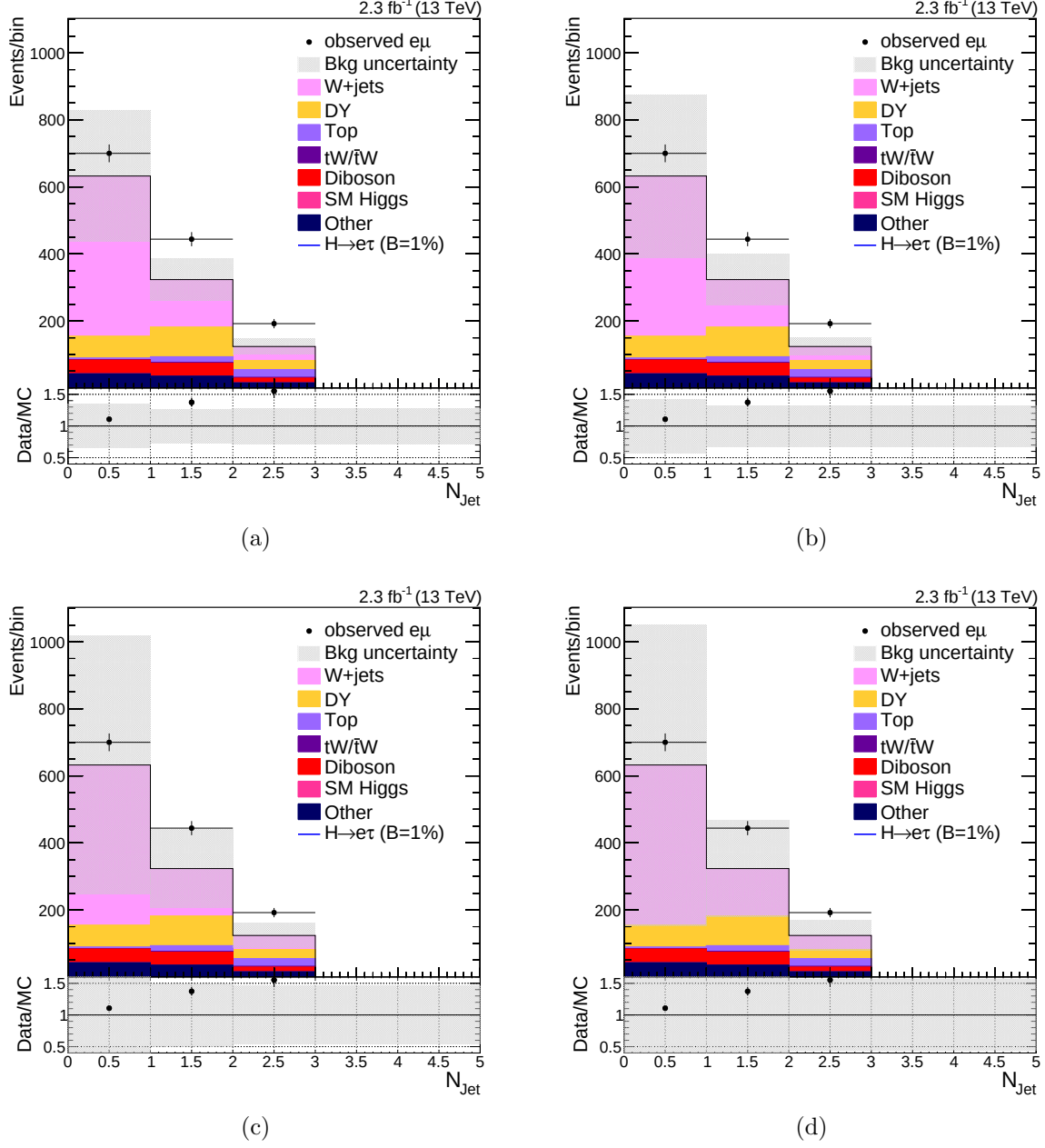


Figure D.5.: Validation of the W+jets normalisation uncertainty. The event yield is shown as function of the jet multiplicity. Background uncertainty is calculated with a normalisation uncertainty for the W+jets background of (a) 40%, (b) 50%, (c) 80%, and (d) 100%.

E. Search for LFV decays of the Higgs boson at $\sqrt{s} = 8$ TeV

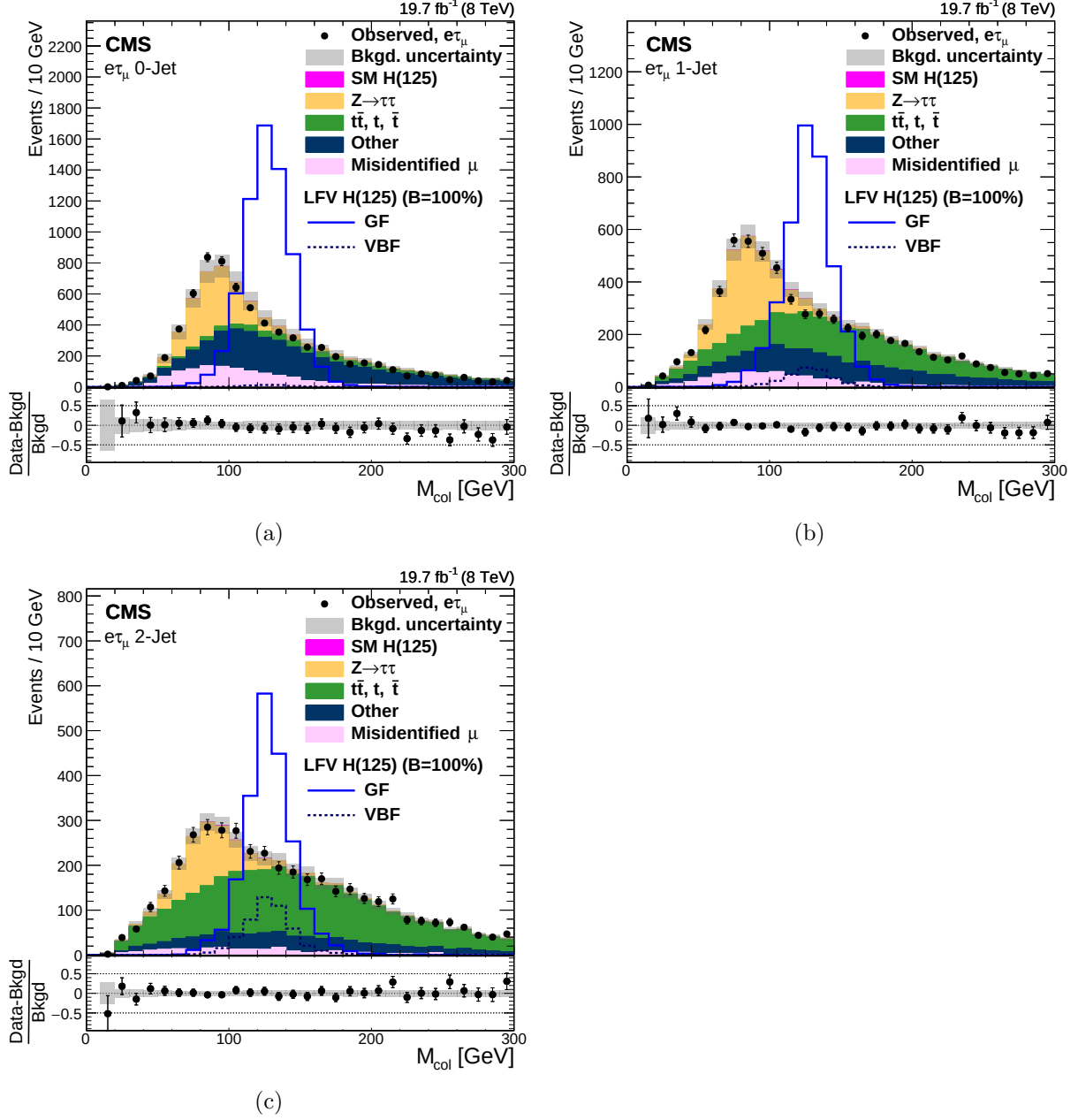


Figure E.1.: Collinear mass distributions after loose selection for the 0-jet category (a), 1-jet category (b), and 2-jet category (c). Observed distributions are compared with the background expectation. The simulated distributions for the signal are shown for the branching fraction $B(H \rightarrow e\tau)=100\%$. Collision data was collected in proton-proton collisions at $\sqrt{s} = 8$ TeV [5].

F. Discriminating variables of the 8 TeV analysis at $\sqrt{s} = 13$ TeV

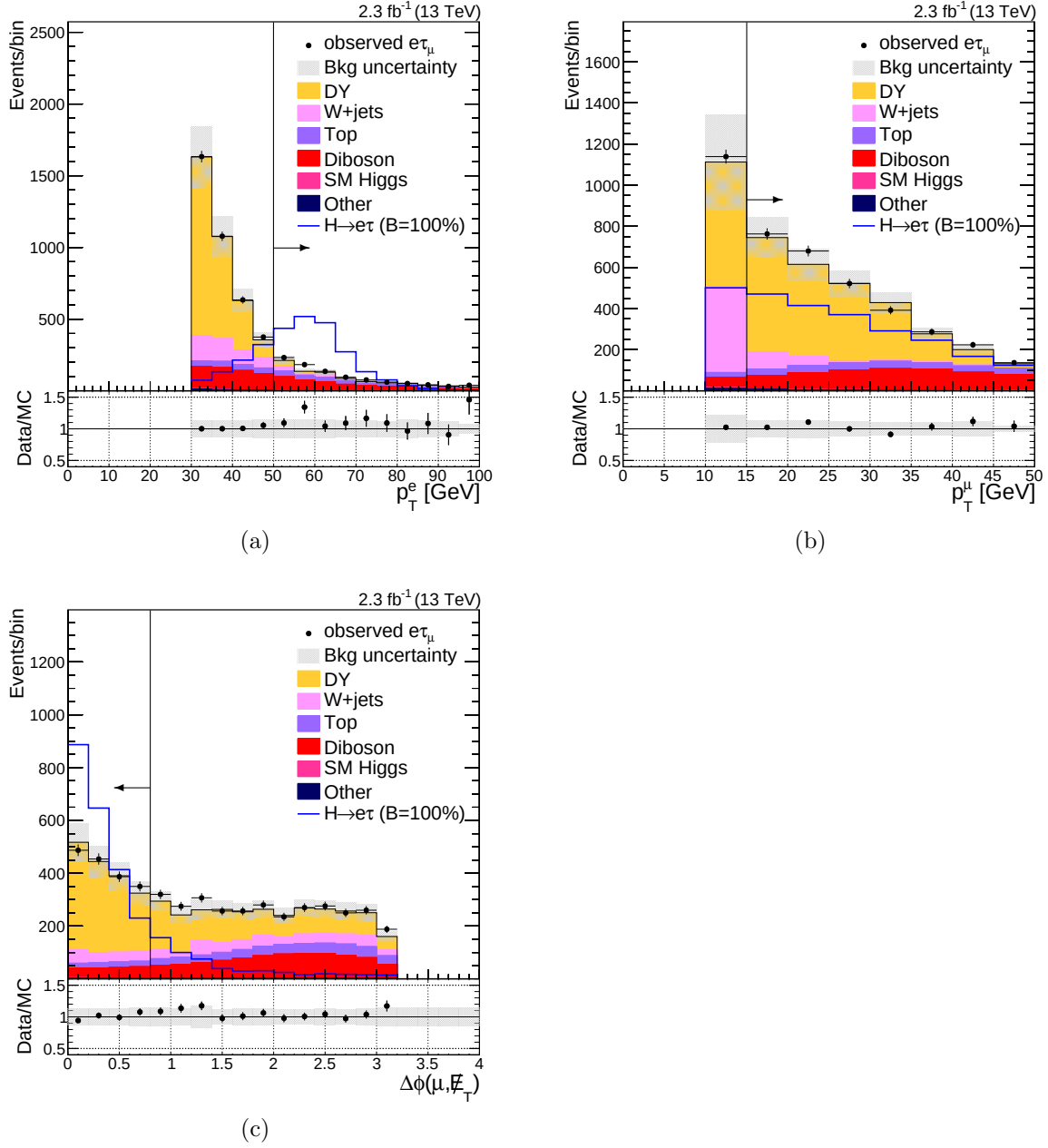


Figure F.1.: Event yields in the 0-jet category after preselection as function of p_T^e (a), p_T^μ (b), and $\Delta\phi(\mu, \cancel{E}_T)$ (c). These variables are the discriminating variables which were used to define the criteria of the event selection for the 0-jet category of the $\sqrt{s} = 8$ TeV analysis. The corresponding criteria are indicated as a solid black line with an arrow indicating the events which are considered in the event selection.

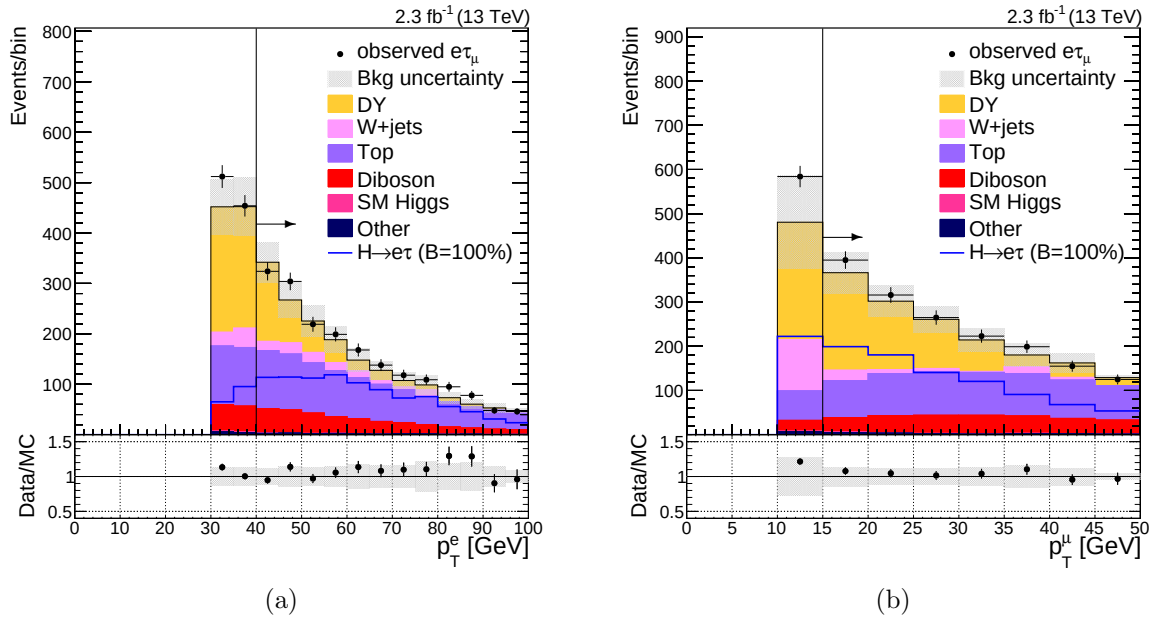


Figure F.2.: Event yields in the 1-jet category after preselection as function of p_T^e (a), and p_T^μ (b). These variables are the discriminating variables which were used to define the criteria of the event selection for the 1-jet category of the $\sqrt{s} = 8$ TeV analysis. The corresponding criteria are indicated as a solid black line with an arrow indicating the events which are considered in the event selection.

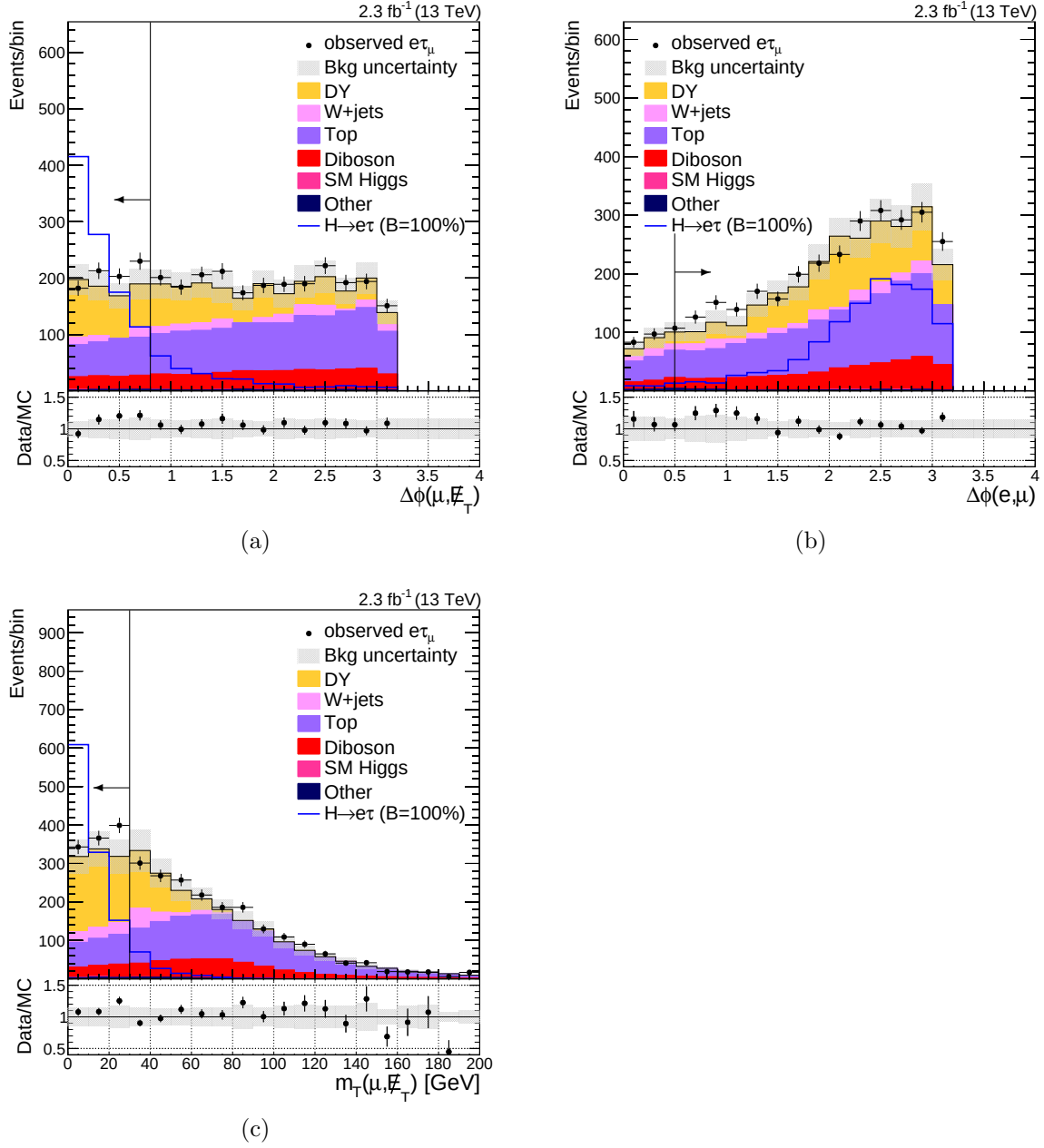


Figure F.3.: Event yields in the 1-jet category after preselection as function of $\Delta\phi(\mu, \cancel{E}_T)$ (a), $\Delta\phi(e, \mu)$ (b), and $m_T(\mu, \cancel{E}_T)$ (c). These variables are the discriminating variables which were used to define the criteria of the event selection for the 1-jet category of the $\sqrt{s} = 8$ TeV analysis. The corresponding criteria are indicated as a solid black line with an arrow indicating the events which are considered in the event selection.

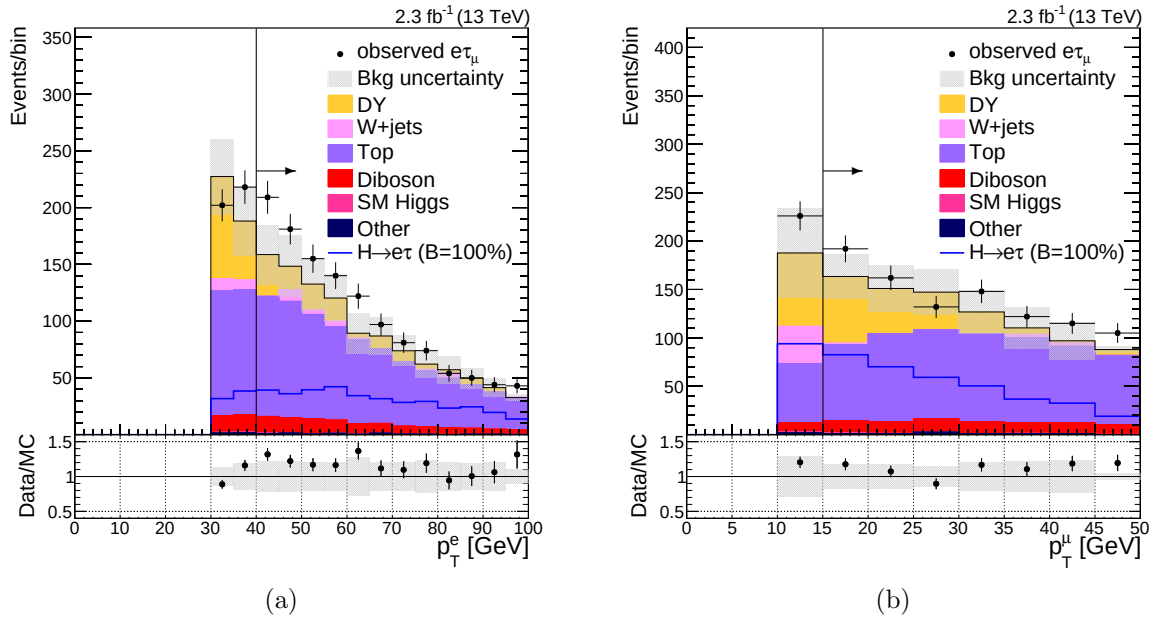


Figure F.4.: Event yields in the 2-jet category after preselection as function of p_T^e (a), and p_T^μ (b). These variables are the discriminating variables which were used to define the criteria of the event selection for the 2-jet category of the $\sqrt{s} = 8$ TeV analysis. The corresponding criteria are indicated as a solid black line with an arrow indicating the events which are considered in the event selection.

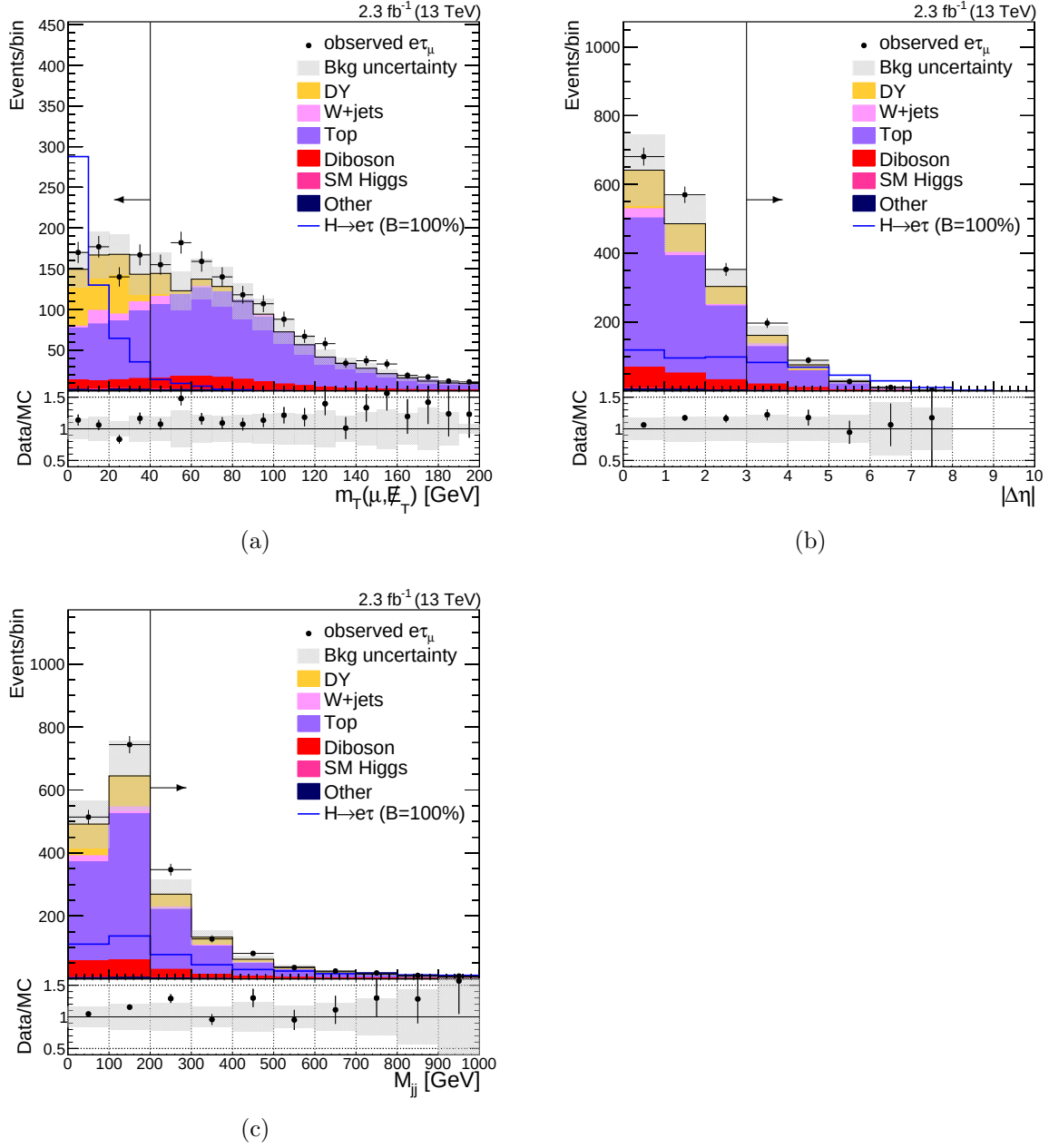


Figure F.5.: Event yields in the 2-jet category after preselection as function of $m_T(\mu, \cancel{E}_T)$ (a), $|\Delta\eta|$ (b), and M_{jj} (c). These variables are the discriminating variables which were used to define the criteria of the event selection for the 2-jet category of the $\sqrt{s} = 8$ TeV analysis. The corresponding criteria are indicated as a solid black line with an arrow indicating the events which are considered in the event selection.

G. Expression for the 95% CL upper limit on the branching fraction

The significance Z is given by:

$$Z = \frac{S}{\sqrt{S+B + \sum_{i=0}^b (\Delta B_i)^2}}, \quad (\text{G.1a})$$

where S is the number of signal events, B is the number of background events, and $(\Delta B_i)^2$ is the systematic uncertainty of background process i .

$$Z^2 = \frac{S^2}{S+B + \sum_{i=0}^b (\Delta B_i)^2} \quad (\text{G.1b})$$

Equation G.1a can be squared and rearranged in order to use the quadratic formula.

$$S^2 - Z^2 \cdot S - Z^2 \cdot \left(B + \sum_{i=0}^b (\Delta B_i)^2 \right) = 0 \quad (\text{G.1c})$$

Then, the number of signal events can be estimated for a given significance and background composition.

$$S = \frac{Z^2}{2} + \sqrt{\frac{Z^4}{4} + Z^2 \cdot \left(B + \sum_{i=0}^b (\Delta B_i)^2 \right)} \quad (\text{G.1d})$$

$$S = \frac{Z^2}{2} + Z \cdot \sqrt{\frac{Z^2}{4} + B + \sum_{i=0}^b (\Delta B_i)^2} \quad (\text{G.1e})$$

The upper limit on the branching fraction with a 95% CL is given by $Z = 1.96$ ([1], Statistics, Table 39.1) and the upper limit can be expressed in per cent by normalising S to the number of signal events for a branching fraction of 1%.

$$L_{B(H \rightarrow e\tau_\mu)}[\%] = \frac{\frac{Z^2}{2} + Z \cdot \sqrt{\frac{Z^2}{4} + B + \sum_{i=0}^b (\Delta B_i)^2}}{S_{B(H \rightarrow e\tau_\mu)=1\%}} \quad (\text{G.1f})$$

H. Results of the search at $\sqrt{s} = 13 \text{ TeV}$

Table H.1.: Event yields in the signal region, $100 \text{ GeV} < M_{\text{col}} < 150 \text{ GeV}$, after the profile likelihood fit of the signal and background contributions. The estimated signal and background contributions were normalised to 2.32 fb^{-1} before the fit. Furthermore, the SM Higgs boson production cross section was assumed, taking gluon fusion and VBF Higgs-production into account. The best fit branching fraction $B(H \rightarrow e\tau) = 0.66\%$ is used for estimating the expected signal yield, which is derived from MC simulation only. The uncertainty on the signal yield does only include uncertainties for the signal model.

Jet category	0-jet	1-jet	2-jet
$Z \rightarrow ll$	19.9 ± 2.6	19.8 ± 3.4	7.4 ± 2.3
Misidentified leptons	0.0 ± 0.0	2.7 ± 2.1	0.0 ± 0.0
$t\bar{t}$, $tW/\bar{t}W$	2.9 ± 0.4	32.0 ± 1.6	12.8 ± 0.6
WW, WZ, ZZ	16.5 ± 0.7	13.1 ± 0.6	2.1 ± 0.2
$W\gamma^{(*)}$	0.2 ± 0.0	0.9 ± 0.2	0.0 ± 0.0
SM H background	0.6 ± 0.0	1.0 ± 0.1	0.5 ± 0.0
Sum of background	40.1 ± 2.8	69.6 ± 4.1	22.9 ± 2.4
Observed	49.0 ± 7.0	67.0 ± 8.2	22.0 ± 4.7
LFV H signal (B=0.66%)	6.6 ± 0.4	4.1 ± 0.2	1.3 ± 0.1

Table H.2.: Comparison of the results for a search for Higgs boson decays to $e\tau_\mu$ at $\sqrt{s} = 8 \text{ TeV}$ and $\sqrt{s} = 13 \text{ TeV}$. Numbers for $\sqrt{s} = 8 \text{ TeV}$ were taken from Table 8 in [5] and they were taken for $\sqrt{s} = 13 \text{ TeV}$ from Table 8.24. The number of signal events (S) is translated into a branching fraction of $B(H \rightarrow e\tau) = 0.66\%$. Fitted event yields are given for the sum of the backgrounds (B), and the $t\bar{t}$ and $tW/\bar{t}W$ background processes (Top).

	$\sqrt{s} = 8 \text{ TeV}$			$\sqrt{s} = 13 \text{ TeV}$		
	0-jet	1-jet	2-jet	0-jet	1-jet	2-jet
Signal	22.1 ± 1.5	15.3 ± 1.1	5.6 ± 1.0	6.6 ± 0.4	4.1 ± 0.2	1.3 ± 0.1
B	273.5 ± 6.1	282.0 ± 6.0	28.1 ± 1.3	42.1 ± 2.1	71.8 ± 3.5	23.3 ± 2.4
S/B [%]	8.08 ± 0.59	5.43 ± 0.42	20.1 ± 3.5	15.7 ± 1.2	5.71 ± 0.39	5.58 ± 0.72
$S/\sqrt{S+B}$	1.29 ± 0.09	0.89 ± 0.07	0.97 ± 0.15	0.95 ± 0.06	0.47 ± 0.02	0.26 ± 0.02
Top ($t\bar{t}$, $tW/\bar{t}W$)	13.8 ± 0.3	69.4 ± 2.3	12.7 ± 0.8	3.0 ± 0.3	32.4 ± 0.9	12.9 ± 0.5
Top/B [%]	5.0 ± 0.2	24.6 ± 1.0	45.2 ± 3.5	7 ± 1	45 ± 3	55 ± 6

Table H.3.: Changes in the signal and background yields between $\sqrt{s} = 13$ TeV and $\sqrt{s} = 8$ TeV analyses given as ratio. Number were taken from Table H.2. S, B, and Top are the event yields of the signal, the sum of backgrounds, and the $t\bar{t}$ and $tW/\bar{t}W$ backgrounds, respectively.

	0-jet	1-jet	2-jet
S	0.3 ± 0.03	0.27 ± 0.02	0.23 ± 0.04
B	0.15 ± 0.01	0.25 ± 0.01	0.83 ± 0.09
S/B	1.9 ± 0.2	1.1 ± 0.2	0.28 ± 0.06
$S/\sqrt{S+B}$	0.74 ± 0.07	0.53 ± 0.05	0.27 ± 0.05
Top/B	1.4 ± 0.2	1.8 ± 0.1	1.2 ± 0.2

Table H.4.: Expected event production at $\sqrt{s} = 8$ TeV and $\sqrt{s} = 13$ TeV. Cross sections for $\sqrt{s} = 8$ TeV ($\sigma_{8\text{TeV}}$) and $\sqrt{s} = 13$ TeV ($\sigma_{13\text{TeV}}$) are listed for the dominant SM Higgs boson production-processes. The expected event yield is calculated with $N = L \cdot \sigma$, where L is the integrated luminosity. The number of expected events for $\sqrt{s} = 8$ TeV ($N_{8\text{TeV}}$) are calculated for the integrated luminosity 19.7 fb^{-1} , while the 2015 data set at $\sqrt{s} = 13$ TeV with an integrated luminosity of 2.32 fb^{-1} is used for the calculation of $N_{13\text{TeV}}$. The ratio of the number of expected events for both run periods is calculated.

Process	$\sigma_{8\text{TeV}}$ [pb]	$\sigma_{13\text{TeV}}$ [pb]	$N_{8\text{TeV}}$	$N_{13\text{TeV}}$	$N_{13\text{TeV}}/N_{8\text{TeV}}$
GF SM Higgs production [142, 145]	19.47	44.14	383,560	102,410	0.27
VBF SM Higgs production [142, 145]	1.601	3.782	31,540	8,774	0.28
$t\bar{t}$ production [129, 146]	245	772	4,826,500	1,791,000	0.37

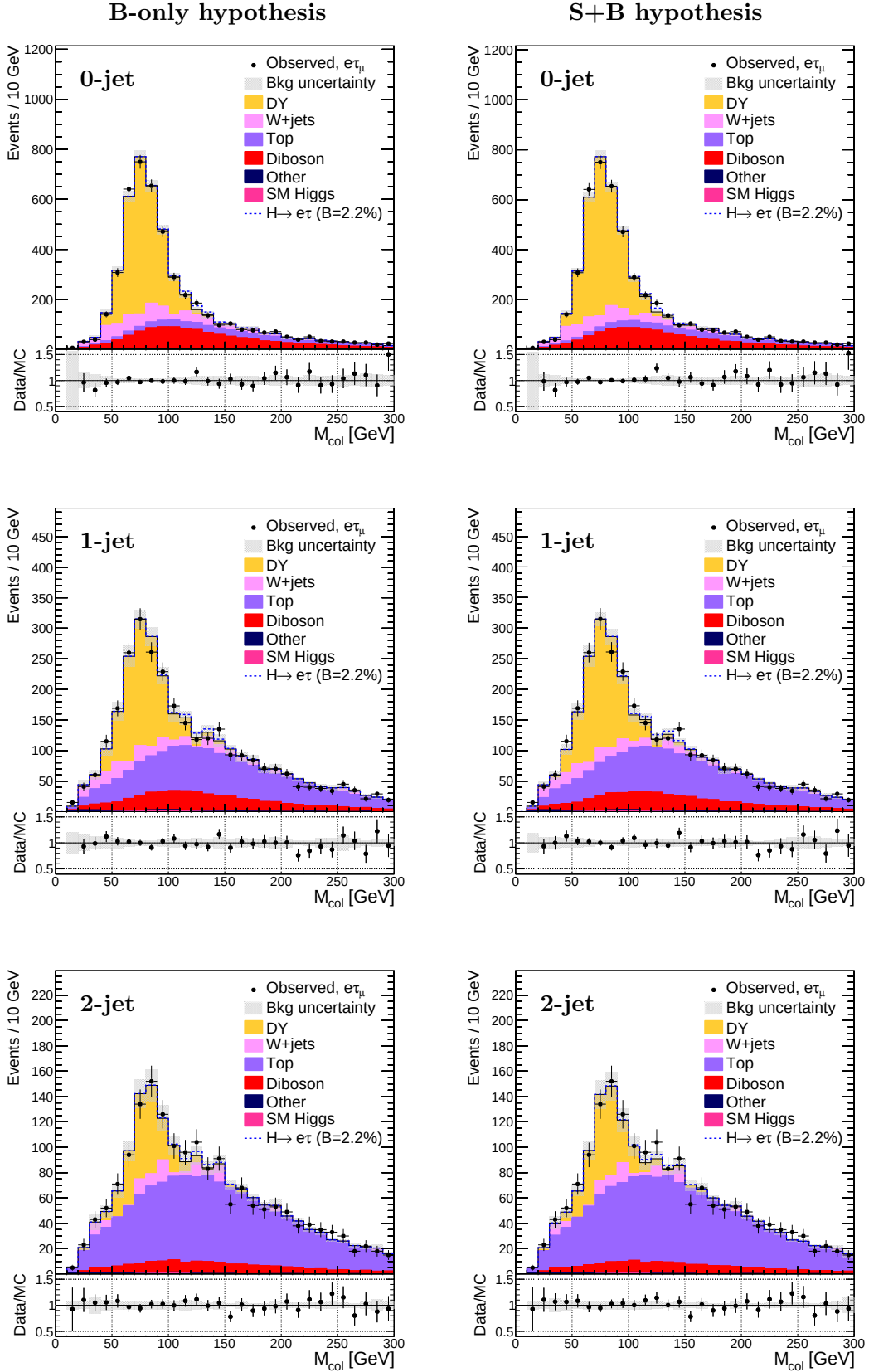


Figure H.1.: Comparison of the observed collinear mass distribution with the background expectations after the binned likelihood fit. Distributions are shown after preselection. The signal expectation is shown for the branching fraction $B(H \rightarrow e\tau) = 0.66\%$. The left column shows the background expectation after the background-only fit, while the right column shows the background expectations after the signal-plus-background fit. In the upper, middle, and lower rows, the collinear mass is shown for the 0-jet, 1-jet, and 2-jet category, respectively.

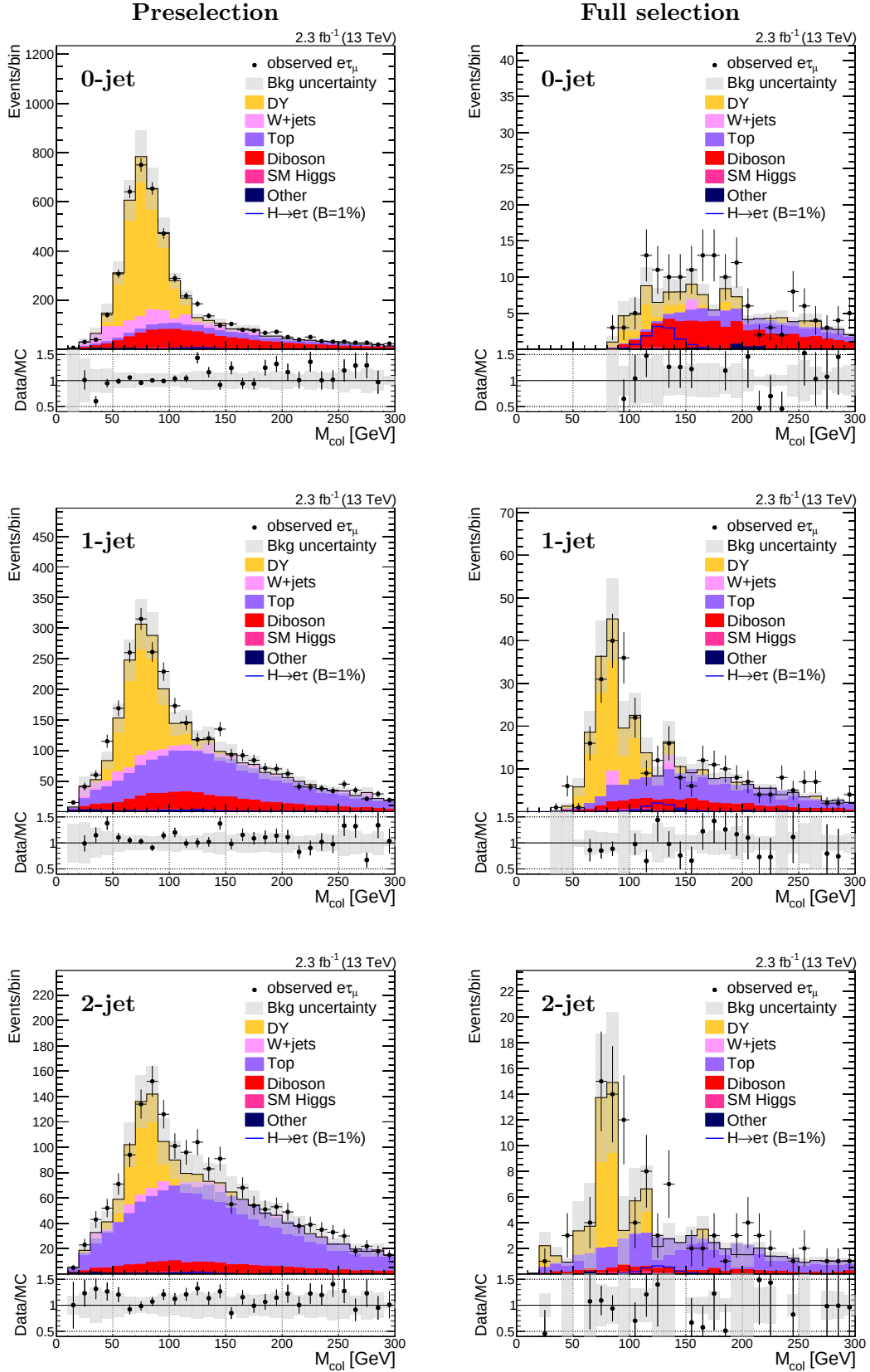


Figure H.2.: Comparison of the observed collinear mass distribution with the background expectations after unblinding. The signal expectation is shown for the branching fraction $B(H \rightarrow e\tau) = 1\%$. The left column shows the background expectation after preselection, while the right column shows the background expectations after applying the full selection. In the upper, middle, and lower rows, the collinear mass is shown for the 0-jet, 1-jet, and 2-jet category, respectively.

I. New luminosity estimate for LHC Run II

Beam-spot movements are measured with the CMS tracking system and corrections to the beam separation are derived. Constraints on the position of the luminous region are made in order to reconstruct the tracks needed in the luminosity measurement of LHC Run II [30]. These constraints were found to be inappropriate, affecting the reconstructed vertex distance scales [30]. A new reconstruction of the data was performed resulting in a 2.7% lower absolute luminosity scale than the original estimate. The luminosity uncertainty of the new measurement was found to be 2.3%, replacing the original uncertainty of 2.7%. This has a small effect on the total integrated luminosity used by the search for LFV decays of the Higgs boson to $e\tau$ in the $e\tau_\mu$ channel. The total integrated luminosity used in the search is 2.26 fb^{-1} instead of 2.32 fb^{-1} .

In the search for LFV decays of the Higgs boson to $e\tau_\mu$, the original luminosity estimate was used for calibration and derivation of scale factors. Cross section measurements using the original luminosity estimate were used to estimate the systematic uncertainty on the normalisation of the background model in control regions. This parts affect mostly differences in collision data and simulation or can be corrected for by the profile-likelihood fit. The optimisation of the event selection remains unaffected as signal and background expectations are both affected. Thus, the effect on the outcome of the optimisation is expected to be small. Only the normalisation of the background models and the nuisance parameter for the luminosity are changed in the fit and the impact on the results of the LFV search is studied.

Table I.1.: The expected upper limits at 95% CL on the branching fraction $B(H \rightarrow e\tau)$ for the different jet categories. The asymmetric one standard-deviation uncertainties around the expected limits are shown in parentheses. These limits were estimated using simulation only. Upper limits on the branching fraction are shown for the original luminosity estimate of 2.32 fb^{-1} and the new estimate of 2.26 fb^{-1} .

Jet category	Expected limit at 95% CL (%)	
	2.32 fb^{-1}	2.26 fb^{-1}
0-jet	$1.46 \left({}^{+0.67}_{-0.44} \right)$	$1.46 \left({}^{+0.68}_{-0.44} \right)$
1-jet	$2.91 \left({}^{+1.38}_{-0.89} \right)$	$2.95 \left({}^{+1.40}_{-0.90} \right)$
2-jet	$4.73 \left({}^{+2.32}_{-1.49} \right)$	$4.80 \left({}^{+2.35}_{-1.51} \right)$
combined	$1.21 \left({}^{+0.55}_{-0.36} \right)$	$1.22 \left({}^{+0.55}_{-0.36} \right)$

I.1. Impact on the search for LFV decays of the Higgs boson

First, the effect on the expected upper limits on the branching fraction $B(H \rightarrow e\tau)$ from simulation only is studied. This gives an estimate of the expected change in the upper limits on the branching fraction for this luminosity correction. Upper limits on the branching fraction of both luminosity estimates are compared in Table I.1. There are small changes in the upper limit and its uncertainty for the single jet categories. The effect on the combined upper limit is negligible.

The profile-likelihood fit is redone using the new luminosity estimate for the normalisation of the background model. There are only small changes with respect to the fit using the original luminosity estimate. The collinear mass distributions are shown after the binned likelihood fit in Figure I.1 for the fit using the luminosity estimate of 2.26 fb^{-1} with an uncertainty of 2.3% for the background-only fits and the signal-plus-background fits in each jet category. These fitted collinear mass distributions are compared to the fitted collinear mass distributions using the luminosity estimate of 2.32 fb^{-1} with an uncertainty of 2.7% for the background-only fit in Figure I.2 and for the signal-plus-background fit in Figure I.3. The changes of the collinear mass are small. The new best fit value is $B(H \rightarrow e\tau) = 0.78 \left({}^{+0.22}_{-0.62} \right)$. It is compatible within uncertainties to the original best fit value $B(H \rightarrow e\tau) = 0.66 \left({}^{+0.38}_{-0.57} \right)$.

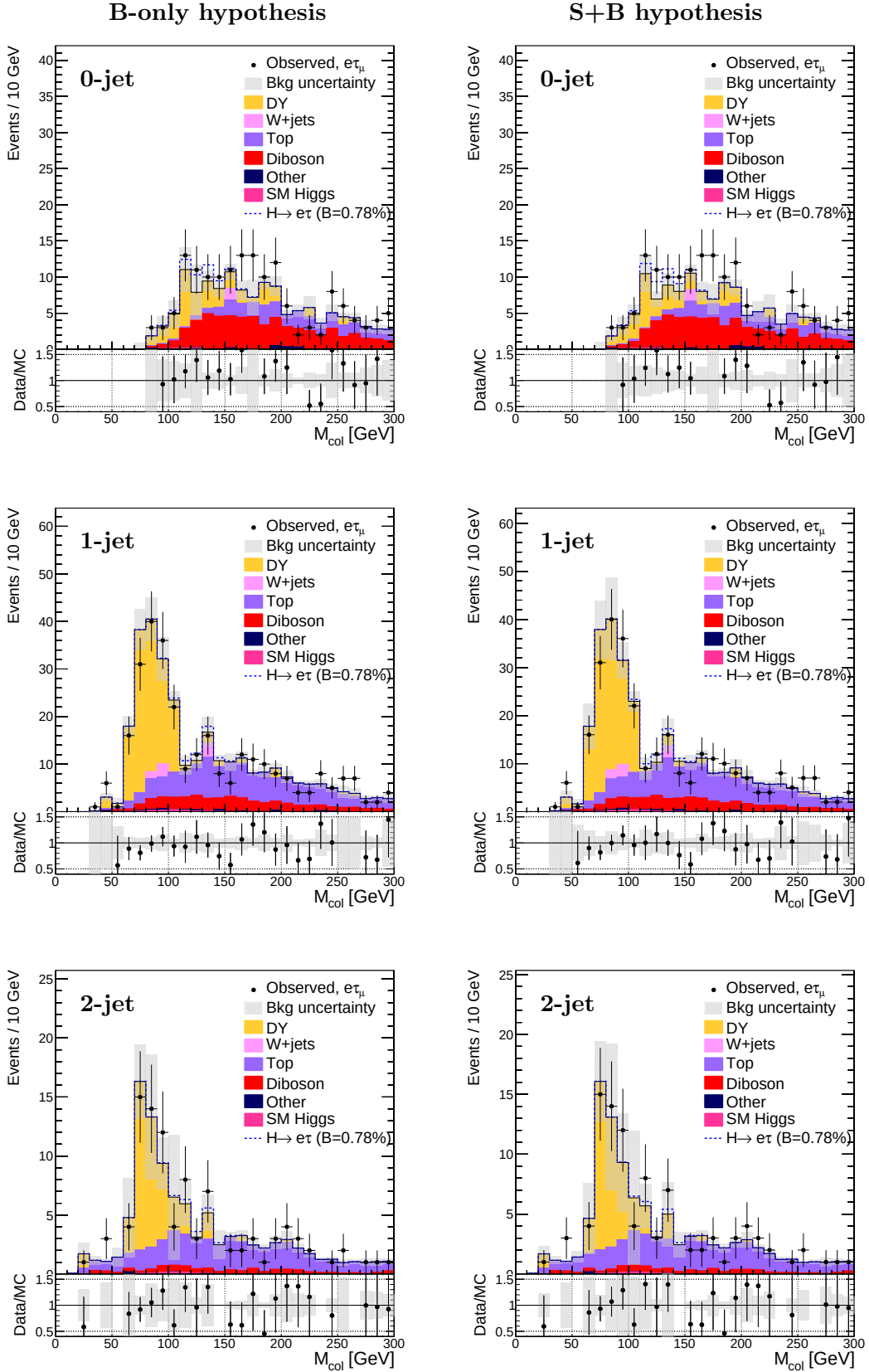


Figure I.1.: Comparison of the observed collinear mass distribution with the background expectations after the binned likelihood fit for a luminosity normalisation of 2.26 fb^{-1} . The signal expectation is shown for the branching fraction $B(H \rightarrow e\tau) = 0.78\%$. The left column shows the background expectation after the background-only fit, while the right column shows the background expectations after the signal-plus-background fit. In the upper, middle, and lower rows, the collinear mass is shown for the 0-jet, 1-jet, and 2-jet category, respectively.

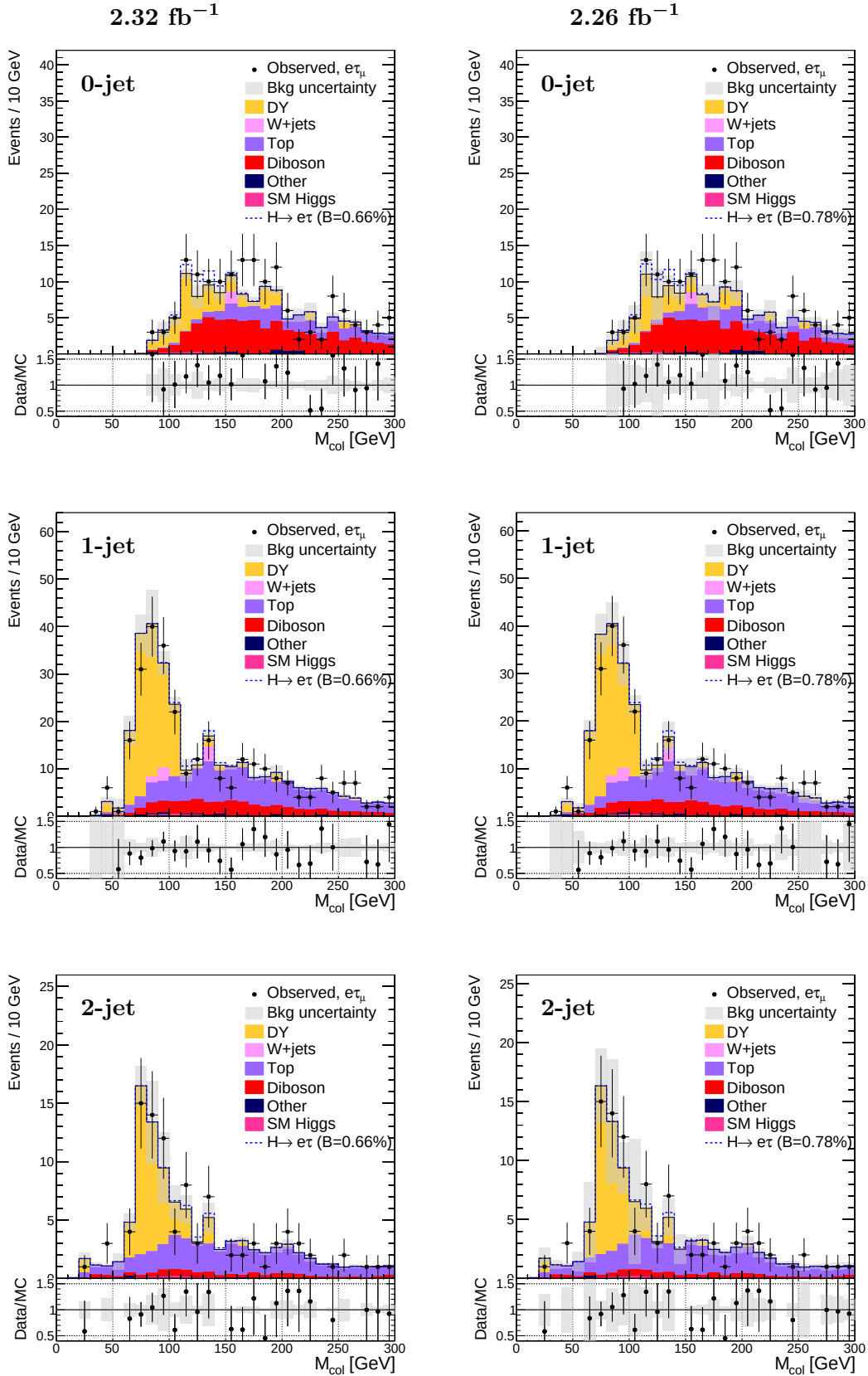


Figure I.2.: Comparison of the observed collinear mass distribution with the background expectations after the binned likelihood fit for a luminosity normalisation of 2.32 fb^{-1} (left column) and 2.26 fb^{-1} (right column), respectively. The signal expectation is shown for the best fit branching fraction $B(H \rightarrow e\tau) = 0.66\%$ and $B(H \rightarrow e\tau) = 0.78\%$, respectively. The background expectation after the background-only fit is shown. In the upper, middle, and lower rows, the collinear mass is shown for the 0-jet, 1-jet, and 2-jet category, respectively.

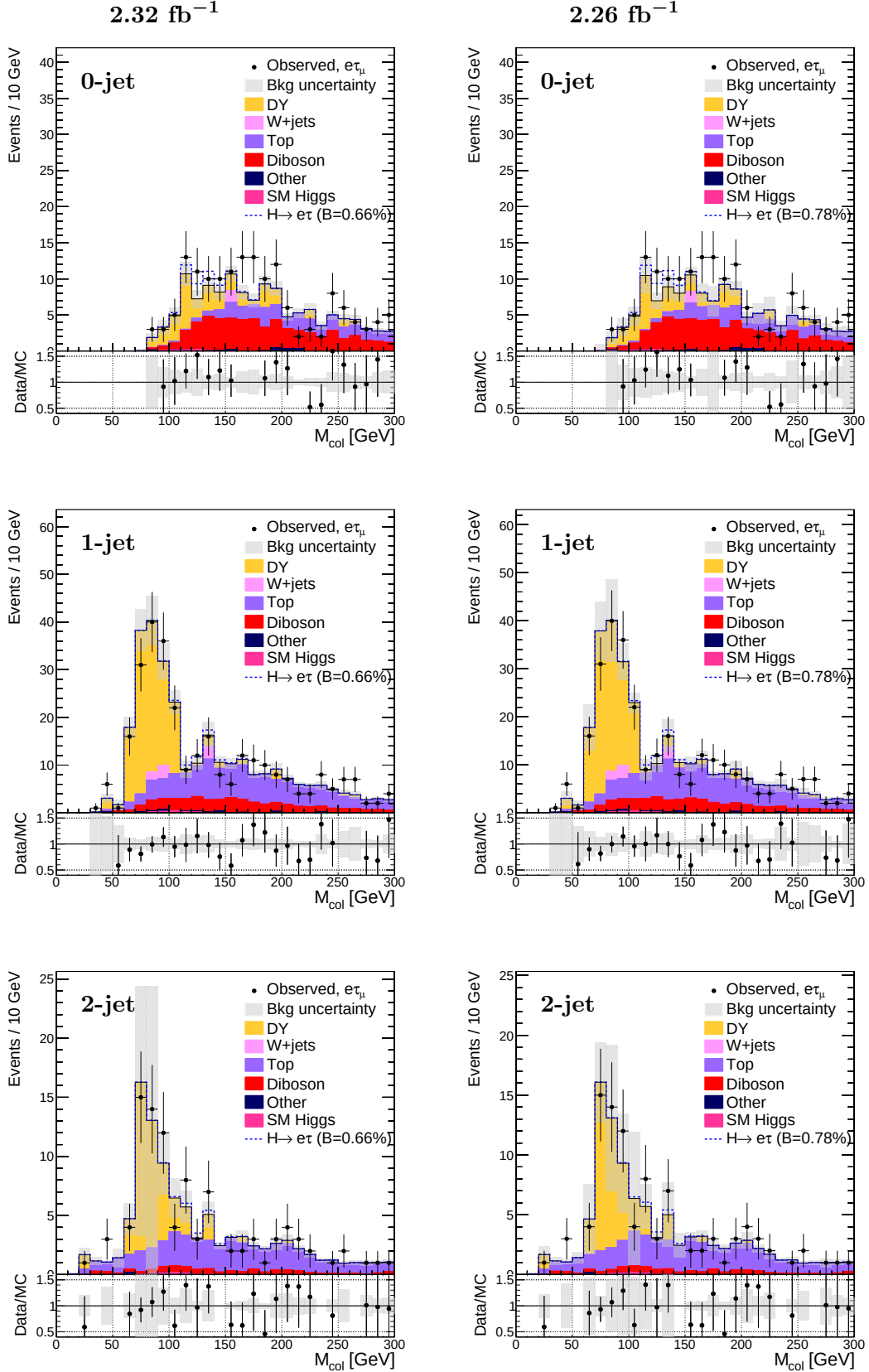


Figure I.3.: Comparison of the observed collinear mass distribution with the background expectations after the binned likelihood fit for a luminosity normalisation of 2.32 fb^{-1} (left column) and 2.26 fb^{-1} (right column), respectively. The signal expectation is shown for the best fit branching fraction $B(H \rightarrow e\tau) = 0.66\%$ and $B(H \rightarrow e\tau) = 0.78\%$, respectively. The background expectation after the signal-plus-background fit is shown. In the upper, middle, and lower rows, the collinear mass is shown for the 0-jet, 1-jet, and 2-jet category, respectively.

I. New luminosity estimate for LHC Run II

Table I.2 lists the event yields in the signal region, $100 \text{ GeV} < M_{\text{col}} < 150 \text{ GeV}$, after the profile-likelihood fit using the original luminosity estimate, while Table I.3 lists the event yield after the profile-likelihood fit using the new luminosity estimate. The background event yield has decreased by less than one event, but the uncertainty has increased.

Table I.2.: Event yields in the signal region, $100 \text{ GeV} < M_{\text{col}} < 150 \text{ GeV}$, after the profile likelihood fit of the background contributions. The estimated background contributions were normalised to 2.32 fb^{-1} before the fit. Furthermore, the SM Higgs boson production cross section was assumed, taking gluon fusion and VBF Higgs-production into account. The obtained best fit branching fraction of the signal-plus-background fit, $B(H \rightarrow e\tau) = 0.66\%$, is used for estimating the expected signal yield.

Jet category	0-jet	1-jet	2-jet
$Z \rightarrow ll$	21.2 ± 2.2	20.8 ± 2.9	7.7 ± 2.2
Misidentified leptons	—	3.2 ± 0.8	—
$t\bar{t}$, $tW/\bar{t}W$	3.0 ± 0.3	32.4 ± 0.9	12.9 ± 0.5
WW , WZ , ZZ	17.1 ± 0.4	13.5 ± 0.6	2.2 ± 0.3
$W\gamma^{(*)}$	0.2 ± 0.0	0.9 ± 0.3	—
SM H background	0.6 ± 0.0	1.0 ± 0.0	0.5 ± 0.0
Sum of background	42.1 ± 2.1	71.8 ± 3.5	23.3 ± 2.4
Observed	49.0 ± 7.0	67.0 ± 8.2	22.0 ± 4.7
LFV $H \rightarrow e\tau$ ($B=0.66\%$)	6.6 ± 0.4	4.1 ± 0.2	1.3 ± 0.1

Table I.3.: Event yields in the signal region, $100 \text{ GeV} < M_{\text{col}} < 150 \text{ GeV}$, after the profile likelihood fit of the background contributions. The estimated background contributions were normalised to 2.26 fb^{-1} before the fit. Furthermore, the SM Higgs boson production cross section was assumed, taking gluon fusion and VBF Higgs-production into account. The obtained best fit branching fraction of the signal-plus-background fit, $B(H \rightarrow e\tau) = 0.78\%$, is used for estimating the expected signal yield.

Jet category	0-jet	1-jet	2-jet
$Z \rightarrow ll$	21.1 ± 5.4	20.8 ± 2.9	7.7 ± 5.9
Misidentified leptons	—	3.0 ± 1.4	—
$t\bar{t}$, $tW/\bar{t}W$	3.0 ± 0.8	32.4 ± 3.1	13.0 ± 0.9
WW , WZ , ZZ	16.9 ± 0.9	13.3 ± 0.8	2.1 ± 0.2
$W\gamma^{(*)}$	0.2 ± 0.0	0.9 ± 0.2	—
SM H background	0.6 ± 0.1	1.0 ± 0.1	0.5 ± 0.0
Sum of background	41.7 ± 6.2	71.4 ± 3.6	23.3 ± 5.8
Observed	49.0 ± 7.0	67.0 ± 8.2	22.0 ± 4.7
LFV $H \rightarrow e\tau$ ($B=0.78\%$)	7.7 ± 0.4	4.8 ± 0.2	1.5 ± 0.1

Table I.4.: The expected and observed upper limits at 95% CL on the branching fraction $B(H \rightarrow e\tau)$ for a Higgs boson with $M_H = 125$ GeV. The limits for each jet categories and the combined limit are shown. Asymmetric one standard-deviation uncertainties on the expected limits are shown in parentheses.

Luminosity [fb^{-1}]	Jet category	0-jet	1-jet	2-jet	combined
2.32	Expected limit at 95% CL (%)	$< 1.60 \left({}^{+0.73}_{-0.48} \right)$	$< 3.20 \left({}^{+1.39}_{-0.92} \right)$	$< 5.14 \left({}^{+2.44}_{-1.61} \right)$	$< 1.31 \left({}^{+0.58}_{-0.39} \right)$
	Observed limit at 95% CL (%)	< 2.23	< 3.04	< 5.94	< 1.94
2.26	Expected limit at 95% CL (%)	$< 1.63 \left({}^{+0.74}_{-0.49} \right)$	$< 3.27 \left({}^{+1.42}_{-0.95} \right)$	$< 5.20 \left({}^{+2.51}_{-1.61} \right)$	$< 1.32 \left({}^{+0.59}_{-0.39} \right)$
	Observed limit at 95% CL (%)	< 2.34	< 3.16	< 6.19	< 2.01

As no excess of events is observed, 95% CL upper limits on the branching fraction $B(H \rightarrow e\tau)$ are derived. They are summarised in Table I.4, where also the upper limits derived using the original luminosity estimate are summarised for comparison. The change in the expected limit between both luminosity estimates is small, while it is larger for the observed luminosity. All new upper limits are weaker than the original ones. The new combined observed (expected) upper limit on $B(H \rightarrow e\tau)$ is 2.01% (1.32%) at 95% CL. Figure I.4 summarises the expected and the observed upper limits at 95% CL for the original and the new luminosity estimate.

I.2. Results

The new luminosity estimate for LHC Run II [30] has only a small effect on the search for lepton-flavour violating decays of the Higgs boson to $e\tau$ in the $e\tau_\mu$ channel. Collision data with a total integrated luminosity of 2.26 fb^{-1} was used in the search with an uncertainty of 2.3% on the luminosity estimate. Only the normalisation of the background model and the nuisance parameter was changed and the profile-likelihood fit was redone. Expected upper limits on the branching fractions were hardly changed, while the new observed upper limits are a bit weaker than the original ones. The new best fit value is $B(H \rightarrow e\tau) = 0.78 \left({}^{+0.22}_{-0.62} \right)$. No excess of event is observed and the new combined observed (expected) upper limit on $B(H \rightarrow e\tau)$ is 2.01% (1.32%) at 95% CL.

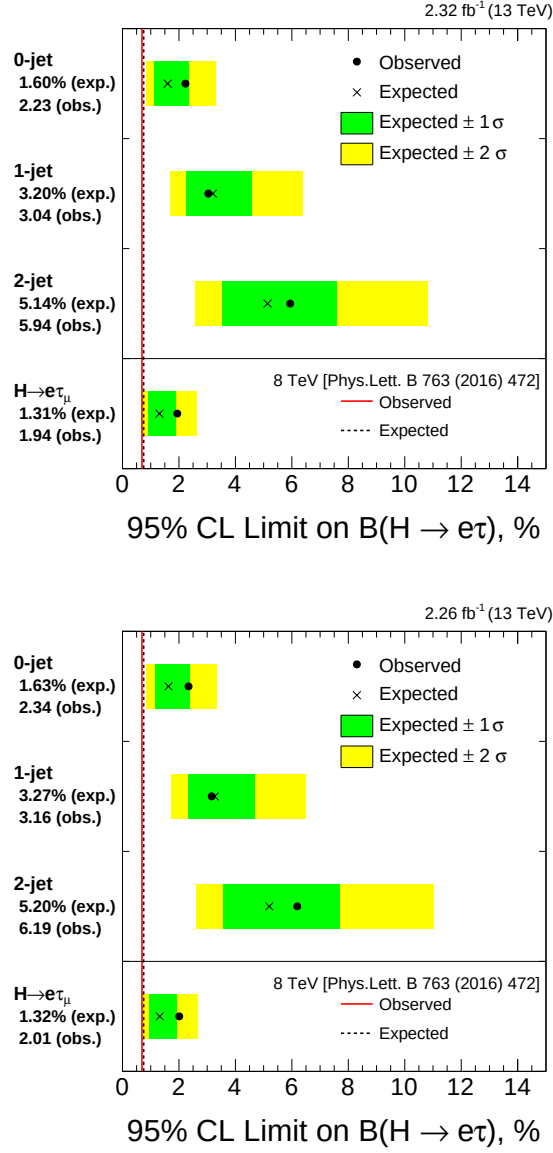


Figure I.4.: The expected and observed upper limits at 95% CL on the branching fraction $B(H \rightarrow e\tau)$ for a Higgs boson with $M_H = 125$ GeV. The limits for each jet categories and the combined limit are shown. Asymmetric one standard-deviation uncertainties on the expected limits are shown in green, while two standard-deviation uncertainties are shown in yellow. The results of the search at $\sqrt{s} = 8$ TeV [5] are indicated as red line (observed combined limit) and black dashed line (expected combined limit). A luminosity estimate of 2.32 fb^{-1} was used for the deriving the upper plot, while a new luminosity estimate of 2.26 fb^{-1} was used to derive the lower plot.

J. Search for LFV decays of the Higgs boson with the ATLAS experiment

The ATLAS collaboration has published direct searches for lepton-flavour violating decays of the Higgs boson to $e\tau$ and $\mu\tau$ [147, 148]. A data sample with an integrated luminosity of 20.3 fb^{-1} was collected in proton-proton collisions at $\sqrt{s} = 8 \text{ TeV}$. Each search was performed in a leptonic and a hadronic channel. The searches in the hadronic channels are model independent, while the leptonic channels [148] are based on the assumption that either the Yukawa couplings for the decay of the Higgs boson to $e\tau$ or the one for the decay to $\mu\tau$ is zero

Search in the hadronic channels The searches in the hadronic channels are model independent. They have two signal regions SR1 and SR2 classified by requirements on $m_T(l, \cancel{E}_T)$ and $m_T(\tau_h, \cancel{E}_T)$, where l is the prompt lepton and $\cancel{E}_T$ the missing energy in the transverse plane. The mass of the Higgs boson is reconstructed using the *Missing Mass Calculator* (MMC) [149]. Shape correction factors for the shape of the W+jets background with respect to the mass are derived from a comparison to a control region defined by requirements on $m_T(l, \cancel{E}_T)$ and $m_T(\tau_h, \cancel{E}_T)$. These corrections depend on the number of jets, the transverse momentum of the tau, and the difference in η between the prompt lepton l and the tau. These shape corrections are applied only in SR1 because they are not needed in SR2. In Section 8.5.2, it is described that there is a dependence of the source for the fake muon on the jet multiplicity, the collinear mass, and $\Delta\phi(e, \mu)$. The scale factor needed for SR1 could be explained by differences in the composition of fake taus in each region, too. It should be studied if the composition of fake taus depends on the MMC mass distribution. The taus of the Higgs boson decays in their signal model were treated as unpolarised because "the left- and right-handed tau polarisation states are produced at equal rates" [147, 148]. This statement is true for the SM but it is not true in the general case. Left- and right-handed polarisation states are produced at equal rates for $H \rightarrow \tau\tau$ because one tau is always left-handed, while the other one is right-handed. However, there are two Yukawa couplings for the LFV decays of the Higgs boson with swapped handedness of the leptons. Both Yukawa couplings do not have to be the same. Though, the effect of polarised and unpolarised taus is negligible for the discussed analyses,

but it can be used for extracting the amplitude of each Yukawa coupling if a signal is found. No significant excess was found and observed upper limits of $B(H \rightarrow \mu\tau) < 1.85$ ($1.24^{+0.50}_{-0.35}$ expected) and $B(H \rightarrow e\tau) < 1.81$ ($2.07^{+0.82}_{-0.58}$ expected) were set.

Search in the leptonic channels The ATLAS searches in the leptonic channels [148] are based on the assumption that either the Yukawa couplings for the decay of the Higgs boson to $e\tau$ or the one for the decay to $\mu\tau$ is zero. A data-driven background method, described in [150], is used for the prompt lepton backgrounds by using the approximation that the corresponding SM background processes are symmetric under replacement of the prompt electron with a prompt muon and vice versa. The prompt lepton background in each leptonic channel is estimated by the prompt lepton yield in the other channel. Corrections for differences in efficiency are applied and the misidentified lepton background is estimated with another method which relies on the difference in identification efficiency between prompt and non-prompt leptons. This method is sensitive to the difference of $B(H \rightarrow e\tau)$ and $B(H \rightarrow \mu\tau)$ and assumes that either one of them is zero to extract an upper limit on the branching fraction. The search is done under the statement [147, 148]: "The bound on $\mu \rightarrow e\gamma$ decays suggest that the presence of a $H \rightarrow \mu\tau$ signal would exclude the presence of a $H \rightarrow e\tau$ signal, and vice versa, at an experimentally observable level at the LHC [16]." However, it is stated in [16] that the bound on $\mu \rightarrow e\gamma$ implies that only $B(H \rightarrow e\tau)$ or $B(H \rightarrow \mu\tau)$ is $\mathcal{O}(10\%)$ and it is not stated that only one of them is non zero at an experimentally observable level at the LHC. Furthermore, the indirect limits on the branching fraction are derived under some assumptions. As the sensitivity of the searches for LFV decays of the Higgs boson to $e\tau$ and $\mu\tau$ has reached $\mathcal{O}(1\%)$, the assumption that one signal yield of one of the decays is negligible has to be verified again. Model dependent observed upper limit limits of $B(H \rightarrow \mu\tau) < 1.79$ ($1.73^{+0.74}_{-0.49}$ expected) and $B(H \rightarrow e\tau) < 1.36$ ($1.48^{+0.60}_{-0.42}$ expected) are derived.

Combination The ATLAS collaboration states upper limits of the combination of the model-independent hadronic channels and the model-dependent leptonic channels. Model dependent observed upper limit limits of $B(H \rightarrow \mu\tau) < 1.43$ ($1.01^{+0.40}_{-0.29}$ expected) and $B(H \rightarrow e\tau) < 1.04$ ($1.21^{+0.49}_{-0.34}$ expected) are derived.

List of Figures

2	The Standard Model	5
2.1.	Elementary particles of the Standard Model	6
2.2.	Fundamental QED vertex	7
2.3.	Fundamental vertices of the weak interaction	8
2.4.	Fundamental vertices of the strong interaction	9
2.5.	Fundamental vertices of the Higgs boson	10
3	Lepton-flavour violating decays of the Higgs boson	19
3.1.	Diagrams contributing to the LFV decay $\tau \rightarrow \mu\gamma$	25
3.2.	Diagrams contributing to the LFV decay $\tau \rightarrow 3\mu$	26
3.3.	Diagram leading to muonium-antimuonium oscillations	27
3.4.	Diagram contributing to magnetic and electric dipole moments of the muon	28
3.5.	Diagrams contributing to $\mu \rightarrow e$ conversion in nuclei	28
3.6.	Low-energy constraints on lepton-flavour violating Yukawa couplings	30
3.7.	Production cross-sections of the Higgs boson at the LHC	32
3.8.	Feynman diagrams of lepton-flavour violating Higgs-boson decays	33
3.9.	Expected branching fractions of the SM Higgs-boson	34
3.10.	Upper limits on the branching fraction for LFV $H \rightarrow \mu\tau$ decays	36
4	Experimental Setup	37
4.1.	The CERN accelerator complex	38
4.2.	Aerial View of the CERN taken in 2008	39
4.3.	Cumulative luminosity as function of time delivered to CMS	42
4.4.	The Compact Muon Solenoid experiment	43
4.5.	Coordinate system convention of CMS	44
4.6.	The CMS detector	47
4.7.	Inner tracking system of CMS	49
4.8.	Layout of the CMS electromagnetic calorimeter	50
4.9.	Location of the ECAL and HCAL detectors in the CMS detector	51
4.10.	Layout of the CMS muon system	53

4.11. Overview of the trigger and data acquisition system	54
5 Event simulation	59
5.1. Monte Carlo simulation of an event in proton-proton collisions	60
6 Event reconstruction	77
6.1. Particle signatures in the CMS detector	78
6.2. Dielectron invariant mass spectrum	91
6.3. Muon sources in simulation	99
6.4. Invariant mass spectrum of dimuons	103
6.5. Jet energy corrections	112
6.6. Pileup observables after pileup reweighting	114
6.7. CSV discriminator distribution for LHC Run I	125
6.8. Performance of the CSV algorithm at 7 TeV	127
6.9. Scale factors SF_b for identification of b-quark jets at $\sqrt{s} = 8$ TeV	128
6.10. Scale factors SF_{light} for identification of b-quark jets at $\sqrt{s} = 8$ TeV	129
6.11. CSV discriminator distribution for different LHC runs	133
7 Impact of heavy flavour modelling on the misidentification of c-quark jets	135
7.1. Dependency of the performance on the charmed hadron type	138
7.2. Charmed hadron decay modelling in PYTHIA 6 and HERWIG 6	141
7.3. Charmed hadron momentum fraction in PYTHIA 6	142
7.4. Charmed hadron momentum fraction in PYTHIA 6 and HERWIG 6	143
7.5. Discriminator and performance in PYTHIA 6 and HERWIG 6	144
7.6. Discriminator distribution before and after reweighting HERWIG 6	147
7.7. Discriminator-distribution differences after reweighting HERWIG 6	149
7.8. Differences between HERWIG 6 and PYTHIA 6 before and after reweighting	151
7.9. Performance of the CSV algorithm before and after reweighting	151
7.10. Measured and predicted charmed hadron production fractions	153
7.11. Effect of reweighting charmed hadron production fractions to ZEUS values	154
7.12. Measured and simulated charmed hadron charged-daughter multiplicities .	157
7.13. Effect of reweighting PYTHIA 6 and HERWIG 6 to measured decays . . .	159
7.14. Peterson/SLAC fragmentation function for heavy flavour	161
7.15. Measured and simulated $D^0/\bar{D}^0$ meson momentum fraction distribution . .	163
7.16. Effect of reweighting momentum fractions to CLEO values	165
7.17. Effect of reweighting to measured ZEUS, PDG and CLEO values	167
7.18. Remaining differences after reweighting to PYTHIA 6 or to measurements	167

7.19. Effect of reweighting on the misidentification probability	171
7.20. Secondary vertex mass and flight distance of c-quark jets	175
8 Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$	177
8.1. Feynman diagram of LFV $H \rightarrow e\tau_\mu$ decay	178
8.2. Shapes of LFV and SM Higgs variable distributions	180
8.3. Estimation of the neutrino momentum	181
8.4. Mass resolution of the visible and collinear mass for the Higgs boson	183
8.5. Feynman diagrams of background processes	185
8.6. Number of interactions before and after pileup reweighting	193
8.7. Electron scale-factors for reconstruction and identification efficiencies . . .	194
8.8. Electron scale-factors for trigger and isolation efficiencies	195
8.9. Muon scale-factors for identification and isolation efficiencies	196
8.10. Composition of the misidentified lepton background	201
8.11. Misidentified lepton regions	202
8.12. Misidentification ratio f_μ	204
8.13. Collinear mass in the same-charge region	205
8.14. Composition of the misidentified-lepton background in the collinear mass .	209
8.15. Sources for misidentified muons	211
8.16. Drell-Yan control region	214
8.17. Top background control region	216
8.18. WW control-region	218
8.19. WZ control-region	219
8.20. ZZ control-region	220
8.21. Composition of the diboson background	221
8.22. W+jets control region	222
8.23. Collinear mass after preselection	228
8.24. Electron transverse momentum after preselection	230
8.25. Muon transverse momentum after preselection	231
8.26. Distributions of discriminating variables used in the $\sqrt{s} = 8$ TeV analysis .	233
8.27. Discriminating VBF variables of the $\sqrt{s} = 8$ TeV analysis	234
8.28. New discriminating variables	236
8.29. Collinear mass distribution in the same-charge region at $\sqrt{s} = 8$ TeV . . .	238
8.30. Optimisation scan for significance	240
8.31. Post-fit collinear mass distributions at $\sqrt{s} = 8$ TeV	243
8.32. Limits on the branching fraction $B(H \rightarrow e\tau)$ at $\sqrt{s} = 8$ TeV	245
8.33. Constraints on the LFV Yukawa couplings at $\sqrt{s} = 8$ TeV	245
8.34. Example optimisation scans for the expected upper limit	249

8.35. Optimisation scans for significance in the 0-jet category	251
8.36. Optimisation scans for expected upper limit in the 0-jet category	253
8.37. Collinear mass after the event selection optimised for $\sqrt{s} = 8$ TeV	257
8.38. Collinear mass after the event selection optimised for significance	259
8.39. Collinear mass after the event selection optimised for expected limit	261
8.40. Collinear mass after the event selection using new variables	263
8.41. Collinear mass distributions after final event selection	265
8.42. Collinear mass distribution after the fit	268
8.44. Expected and observed upper limits on $B(H \rightarrow e\tau)$	272
8.45. Expected limit for $\sqrt{s} = 13$ TeV as function of the integrated luminosity	274
9 Summary	277
9.1. Upper limits on the branching fraction of LFV Higgs boson decays	281
Appendix	284
A.1. Mass resolution of the visible and collinear mass for $e\tau_\mu$	286
A.2. Mass resolution of the visible and collinear mass for $\tau_{e\tau_\mu}$	287
C.1. Misidentified-lepton background composition in the M_{col} distribution	292
C.2. Muon scale-factors for the loose-isolation requirement	293
C.3. Validation distributions of misidentified-lepton background	294
C.4. Validation of the misidentified-lepton background in the $\cancel{E}_T$ distribution	295
C.5. Validation distributions of the isolation side-band method	296
C.6. Misidentified-lepton background composition in the N_{Jet} distribution	297
C.7. Region comparison of misidentified-muon background distributions	298
D.1. Distribution of p_T^e in the Drell-Yan control region	300
D.2. Validation of the top-quark pair production process	301
D.3. Jet multiplicity in the W^+W^- process control region	301
D.4. Composition of the diboson background as function of collinear mass	302
D.5. Validation of the W +jets normalisation uncertainty	303
E.1. Collinear mass distributions after loose selection at $\sqrt{s} = 8$ TeV	305
F.1. Discriminating variables of the $\sqrt{s} = 8$ TeV 0-jet category selection	307
F.2. Lepton transverse momenta with $\sqrt{s} = 8$ TeV 1-jet category selection	308
F.3. Discriminating variables of the $\sqrt{s} = 8$ TeV 1-jet category selection	309
F.4. Lepton transverse momenta with $\sqrt{s} = 8$ TeV 2-jet category selection	310
F.5. Discriminating variables of the $\sqrt{s} = 8$ TeV 2-jet category selection	311

H.1. Collinear mass distribution for the preselection after the fit	317
H.2. Unblinded distribution of the collinear mass	318
Appendix I New luminosity estimate for LHC Run II	319
I.1. Collinear mass distribution after the fit	321
I.2. Collinear mass distribution after the fit	322
I.3. Collinear mass distribution after the fit	323
I.4. Expected and observed upper limits on $B(H \rightarrow e\tau)$	327

List of Tables

2	The Standard Model	5
2.1.	Overview of the symmetry groups of the Standard Model	15
3	Lepton-flavour violating decays of the Higgs boson	19
3.1.	Experimental bounds on $B(\ell_i \rightarrow \ell_j \gamma)$	26
3.2.	Experimental bounds on $B(\ell_i \rightarrow 3\ell_j)$	26
3.3.	Low-energy constraints on lepton-flavour violating Yukawa couplings	29
4	Experimental Setup	37
4.1.	Certified data used in the search for LFV decays of the Higgs boson	42
5	Event simulation	59
5.1.	Monte Carlo generators	60
6	Event reconstruction	77
6.1.	Threshold values used in the electron-energy clustering-algorithms	86
7	Impact of heavy flavour modelling on the misidentification of c-quark jets	135
7.1.	Charmed hadrons and their properties	136
7.2.	Simulated QCD multijet samples for heavy flavour modelling study	137
7.3.	Charmed hadron type fractions in PYTHIA 6 and HERWIG 6	140
7.4.	Misidentification probability of c-quark jets for each CSV working point . .	145
7.5.	Measured and predicted charmed hadron production fractions	153
7.6.	Measured charged-daughter multiplicity of $D^0/\bar{D}^0$ mesons	155
7.7.	Measured charged-daughter multiplicity used for reweighting	156
7.8.	Maximum positions of the Peterson fragmentation function	162
7.9.	Effects on the misidentification probability reweighting PYTHIA 6	169
7.10.	Effects on the misidentification probability reweighting HERWIG 6	170
7.11.	Reweighting results for PYTHIA 6	173
7.12.	Reweighting results for HERWIG 6	173

7.13. Differences in the jet flavour fractions	176
8 Search for lepton-flavour violating decays of the Higgs boson to $e\tau_\mu$	177
8.1. Overview of properties of the background processes	186
8.2. Triggering MVA selection for 2015 data and 25 ns	189
8.3. MVA discriminator requirements for the triggering MVA ID wp 80	189
8.4. Medium Muon ID	190
8.5. Loose Jet ID	191
8.6. Overview signal region and control regions	213
8.7. Normalisation uncertainties	224
8.8. Background-normalisation uncertainties	225
8.9. Theory uncertainties of the Higgs boson processes	225
8.10. Event selection of the $\sqrt{s} = 8$ TeV analysis	240
8.13. The expected and observed upper limits on $B(H \rightarrow e\tau)$ at $\sqrt{s} = 8$ TeV	242
8.14. Ranking of the optimisation for significance in the 0-jet category	250
8.15. Event selection optimised for maximising significance	252
8.16. Event selection optimised for expected limit on the branching fraction	254
8.17. Event selection optimised using new discriminating variables	254
8.18. Expected limits of the $\sqrt{s} = 8$ TeV event selection at $\sqrt{s} = 13$ TeV	256
8.19. Expected limits of the event selection optimised for significance	258
8.20. Expected limit of the event selection optimised for expected limit	260
8.21. Expected limits of the event selection using new discriminating variables	262
8.22. Final event selection at $\sqrt{s} = 13$ TeV	264
8.23. Expected upper limits on $B(H \rightarrow e\tau)$ at $\sqrt{s} = 13$ TeV	266
8.24. Event yields in the signal region after fitting the background contributions	271
8.25. Expected and observed limits on $B(H \rightarrow e\tau)$ at $\sqrt{s} = 13$ TeV	271
Appendix	284
B.1. Overview of MC samples used in the LFV analysis	290
B.2. Cross sections	291
H.1. Event yields in the signal region after signal-plus-background fit	315
H.2. Comparison of the results at $\sqrt{s} = 8$ TeV and $\sqrt{s} = 13$ TeV	315
H.3. Changes in the signal and background yields between the two analyses	316
H.4. Expected event production at $\sqrt{s} = 8$ TeV and $\sqrt{s} = 13$ TeV	316

Appendix I New luminosity estimate for LHC Run II
319

I.1.	Expected limits on $B(H \rightarrow e\tau)$	320
I.2.	Event yields in the signal region after background-only fit for 2.32 fb^{-1} . .	325
I.3.	Event yields in the signal region after background-only fit for 2.26 fb^{-1} . .	325
I.4.	Expected and observed limits on $B(H \rightarrow e\tau)$	326

Bibliography

- [1] K. A. Olive, “Review of Particle Physics,” *Chin. Phys. C* **40** (2016) no.10, 100001. doi:10.1088/1674-1137/40/10/100001
- [2] M. Raidal *et al.*, “Flavour physics of leptons and dipole moments,” *Eur. Phys. J. C* **57** (2008) 13 doi:10.1140/epjc/s10052-008-0715-2 [arXiv:0801.1826 [hep-ph]].
- [3] G. Aad *et al.* [ATLAS Collaboration], “Observation of a new particle in the search for the Standard Model Higgs boson with the ATLAS detector at the LHC,” *Phys. Lett. B* **716** (2012) 1 doi:10.1016/j.physletb.2012.08.020 [arXiv:1207.7214 [hep-ex]].
- [4] S. Chatrchyan *et al.* [CMS Collaboration], “Observation of a new boson at a mass of 125 GeV with the CMS experiment at the LHC,” *Phys. Lett. B* **716** (2012) 30 doi:10.1016/j.physletb.2012.08.021 [arXiv:1207.7235 [hep-ex]].
- [5] V. Khachatryan *et al.* [CMS Collaboration], “Search for lepton flavour violating decays of the Higgs boson to $e\tau$ and $e\mu$ in proton-proton collisions at $\sqrt{s} = 8$ TeV,” *Phys. Lett. B* **763** (2016) 472 doi:10.1016/j.physletb.2016.09.062 [arXiv:1607.03561 [hep-ex]].
- [6] R. Harnik, J. Kopp and J. Zupan, “Flavor Violating Higgs Decays,” *JHEP* **1303** (2013) 026 doi:10.1007/JHEP03(2013)026 [arXiv:1209.1397 [hep-ph]].
- [7] L. Evans and P. Bryant, “LHC Machine,” *JINST* **3** (2008) S08001. doi:10.1088/1748-0221/3/08/S08001
- [8] Annika Vanhoefer, ”Search for Supersymmetry at the LHC with Multivariate Analysis Techniques“, master thesis, 2013
- [9] David Griffiths, ”Introduction to Elementary Particles”, Wiley-VCH, Second Edition 2008. ISBN 978-3-527-40601-2
- [10] Matthew Robinson, ”Symmetry and the Standard Model“, Springer, 2011, ISBN 978-1-4419-8266-7
- [11] J. Goldstone, “Field Theories with Superconductor Solutions,” *Nuovo Cim.* **19** (1961) 154. doi:10.1007/BF02812722

- [12] J. Goldstone, A. Salam and S. Weinberg, “Broken Symmetries,” *Phys. Rev.* **127** (1962) 965. doi:10.1103/PhysRev.127.965
- [13] S. Chatrchyan *et al.* [CMS Collaboration], “Observation of a new boson with mass near 125 GeV in pp collisions at $\sqrt{s} = 7$ and 8 TeV,” *JHEP* **1306** (2013) 081 doi:10.1007/JHEP06(2013)081 [arXiv:1303.4571 [hep-ex]].
- [14] P. A. R. Ade *et al.* [Planck Collaboration], “Planck 2015 results. XIII. Cosmological parameters,” *Astron. Astrophys.* **594** (2016) A13 doi:10.1051/0004-6361/201525830 [arXiv:1502.01589 [astro-ph.CO]].
- [15] http://www.nobelprize.org/nobel_prizes/physics/laureates/2011/ (09.01.2017)
- [16] G. Blankenburg, J. Ellis and G. Isidori, “Flavour-Changing Decays of a 125 GeV Higgs-like Particle,” *Phys. Lett. B* **712** (2012) 386 doi:10.1016/j.physletb.2012.05.007 [arXiv:1202.5704 [hep-ph]].
- [17] V. Khachatryan *et al.* [CMS Collaboration], “Search for Lepton-Flavour-Violating Decays of the Higgs Boson,” *Phys. Lett. B* **749** (2015) 337 doi:10.1016/j.physletb.2015.07.053 [arXiv:1502.07400 [hep-ex]].
- [18] J. L. Diaz-Cruz and J. J. Toscano, “Lepton flavor violating decays of Higgs bosons beyond the standard model,” *Phys. Rev. D* **62** (2000) 116005 doi:10.1103/PhysRevD.62.116005 [hep-ph/9910233].
- [19] T. P. Cheng and M. Sher, “Mass Matrix Ansatz and Flavor Nonconservation in Models with Multiple Higgs Doublets,” *Phys. Rev. D* **35** (1987) 3484. doi:10.1103/PhysRevD.35.3484
- [20] K. Nakamura *et al.* [Particle Data Group], “Review of particle physics,” *J. Phys. G* **37** (2010) 075021. doi:10.1088/0954-3899/37/7A/075021
- [21] J. Adam *et al.* [MEG Collaboration], “New limit on the lepton-flavour violating decay $\mu^+ \rightarrow e^+ \gamma$,” *Phys. Rev. Lett.* **107** (2011) 171801 doi:10.1103/PhysRevLett.107.171801 [arXiv:1107.5547 [hep-ex]].
- [22] T. Goto, R. Kitano and S. Mori, “Lepton flavor violating Z -boson couplings from nonstandard Higgs interactions,” *Phys. Rev. D* **92** (2015) 075021 doi:10.1103/PhysRevD.92.075021 [arXiv:1507.03234 [hep-ph]].
- [23] U. Bellgardt *et al.* [SINDRUM Collaboration], “Search for the Decay $\mu^+ \rightarrow e^+ e^+ e^-$,” *Nucl. Phys. B* **299** (1988) 1. doi:10.1016/0550-3213(88)90462-2

- [24] K. Hayasaka *et al.*, “Search for Lepton Flavor Violating Tau Decays into Three Leptons with 719 Million Produced Tau+Tau- Pairs,” *Phys. Lett. B* **687** (2010) 139 doi:10.1016/j.physletb.2010.03.037 [arXiv:1001.3221 [hep-ex]].
- [25] L. Willmann *et al.*, “New bounds from searching for muonium to anti-muonium conversion,” *Phys. Rev. Lett.* **82** (1999) 49 doi:10.1103/PhysRevLett.82.49 [hep-ex/9807011].
- [26] <https://twiki.cern.ch/twiki/bin/view/LHCPhysics/LHCHXSWG> (7 October 2016)
- [27] D. de Florian *et al.* [LHC Higgs Cross Section Working Group], “Handbook of LHC Higgs Cross Sections: 4. Deciphering the Nature of the Higgs Sector,” arXiv:1610.07922 [hep-ph].
- [28] CMS Collaboration, “The CMS experiment at the CERN LHC” *JINST* **3** (2008) S08004. doi:10.1088/1748-0221/3/08/S08004.
- [29] CMS Collaboration, “CMS Detector Performance and Software”, volume 1 of CMS Physics Technical Design Report. CERN, Geneva, 2006. CMS-TDR-008-1. CERN-LHCC-2006-001.
- [30] CMS Collaboration [CMS Collaboration], “CMS Luminosity Measurement for the 2015 Data Taking Period,” CMS-PAS-LUM-15-001.¹
- [31] CMS Collaboration [CMS Collaboration], “CMS Luminosity Based on Pixel Cluster Counting - Summer 2013 Update,” CMS-PAS-LUM-13-001.
- [32] CMS Collaboration [CMS Collaboration], “CMS Luminosity Based on Pixel Cluster Counting - Summer 2012 Update,” CMS-PAS-LUM-12-001.
- [33] S. van der Meer, “Calibration of the Effective Beam Height in the ISR,” CERN-ISR-PO-68-31.
- [34] C. De Melis, “The CERN accelerator complex“, OPEN-PHO-ACCEL-2016-009, 2016
- [35] “The Large Electron-Positron Collider”: <http://cds.cern.ch/record/1997351>, July 2012
- [36] “Linear accelerator 2”: <http://cds.cern.ch/record/1997427>, September 2012
- [37] <https://home.cern/about/accelerators> (09.02.2017)
- [38] “The Proton Synchrotron Booster”: <http://cds.cern.ch/record/1997372>, July 2012
- [39] M. Brice, “Aerial View of the CERN taken in 2008“, CERN-MI-0807031, 2008

¹The CMS documentation [30] was replaced by a new version in February 2017 [30].

- [40] “The Proton Synchrotron”: <http://cds.cern.ch/record/1997189>, January 2012
- [41] “The Super Proton Synchrotron”: <http://cds.cern.ch/record/1997188>, January 2012
- [42] G. Aad *et al.* [ATLAS Collaboration], “The ATLAS Experiment at the CERN Large Hadron Collider,” JINST **3** (2008) S08003. doi:10.1088/1748-0221/3/08/S08003
- [43] K. Aamodt *et al.* [ALICE Collaboration], “The ALICE experiment at the CERN LHC,” JINST **3** (2008) S08002. doi:10.1088/1748-0221/3/08/S08002
- [44] A. A. Alves, Jr. *et al.* [LHCb Collaboration], “The LHCb Detector at the LHC,” JINST **3** (2008) S08005. doi:10.1088/1748-0221/3/08/S08005
- [45] G. Anelli *et al.* [TOTEM Collaboration], “The TOTEM experiment at the CERN Large Hadron Collider,” JINST **3** (2008) S08007. doi:10.1088/1748-0221/3/08/S08007
- [46] O. Adriani *et al.* [LHCf Collaboration], “The LHCf detector at the CERN Large Hadron Collider,” JINST **3** (2008) S08006. doi:10.1088/1748-0221/3/08/S08006
- [47] B. Goddard, W. Herr, E. Keil, M. Lamont, M. Meddahi and E. Peschardt, “Bunch trains for LEP,” Part. Accel. **57** (1997) 237.
- [48] <https://twiki.cern.ch/twiki/bin/view/CMSPublic/LumiPublicResults> (15.03.2017)
- [49] M. Hoch, M. Brice, “CMS experiment“, CMS-PHO-GEN-2008-028, 2008
- [50] V. Khachatryan *et al.* [CMS Collaboration], “Performance of Photon Reconstruction and Identification with the CMS Detector in Proton-Proton Collisions at $\sqrt{s} = 8$ TeV,” JINST **10** (2015) no.08, P08010 doi:10.1088/1748-0221/10/08/P08010 [arXiv:1502.02702 [physics.ins-det]].
- [51] <http://cms.web.cern.ch/news/cms-detector-design> (01.02.2017)
- [52] CERN Publications, “CMS ready for winding up,” BUL-NA-2003-150 (2003)
- [53] V. Khachatryan *et al.* [CMS Collaboration], “CMS Tracking Performance Results from early LHC Operation,” Eur. Phys. J. C **70** (2010) 1165 doi:10.1140/epjc/s10052-010-1491-3 [arXiv:1007.1988 [physics.ins-det]].
- [54] [CMS Collaboration], “Electron reconstruction and identification at $\sqrt{s} = 7$ TeV,” CMS-PAS-EGM-10-004.
- [55] V. Khachatryan *et al.* [CMS Collaboration], “Performance of Electron Reconstruction and Selection with the CMS Detector in Proton-Proton Collisions at $\sqrt{s} = 8$ TeV,” JINST **10** (2015) no.06, P06005 doi:10.1088/1748-0221/10/06/P06005 [arXiv:1502.02701 [physics.ins-det]].

- [56] S. Abdullin *et al.* [USCMS and ECAL/HCAL Collaborations], “The CMS barrel calorimeter response to particle beams from 2-GeV/c to 350-GeV/c,” *Eur. Phys. J. C* **60** (2009) 359 Erratum: [*Eur. Phys. J. C* **61** (2009) 353]. doi:10.1140/epjc/s10052-009-0959-5, 10.1140/epjc/s10052-009-1024-0
- [57] T. Sjostrand, S. Mrenna and P. Z. Skands, “PYTHIA 6.4 Physics and Manual,” *JHEP* **0605** (2006) 026 doi:10.1088/1126-6708/2006/05/026 [hep-ph/0603175].
- [58] G. Corcella, I. G. Knowles, G. Marchesini, S. Moretti, K. Odagiri, P. Richardson, M. H. Seymour and B. R. Webber, “HERWIG 6: An Event generator for hadron emission reactions with interfering gluons (including supersymmetric processes),” *JHEP* **0101** (2001) 010 doi:10.1088/1126-6708/2001/01/010 [hep-ph/0011363].
- [59] K. A. Olive *et al.* [Particle Data Group Collaboration], “Review of Particle Physics,” *Chin. Phys. C* **38** (2014) 090001. doi:10.1088/1674-1137/38/9/090001
- [60] J. Alwall *et al.*, “A Standard format for Les Houches event files,” *Comput. Phys. Commun.* **176** (2007) 300 doi:10.1016/j.cpc.2006.11.010 [hep-ph/0609017].
- [61] J. C. Collins, “Sudakov form-factors,” *Adv. Ser. Direct. High Energy Phys.* **5** (1989) 573 doi:10.1142/9789814503266_0006 [hep-ph/0312336].
- [62] G. Altarelli and G. Parisi, “Asymptotic Freedom in Parton Language,” *Nucl. Phys. B* **126** (1977) 298. doi:10.1016/0550-3213(77)90384-4
- [63] S. Hoeche, F. Krauss, N. Lavesson, L. Lonnblad, M. Mangano, A. Schalicke and S. Schumann, “Matching parton showers and matrix elements,” hep-ph/0602031.
- [64] C. Peterson, D. Schlatter, I. Schmitt and P. M. Zerwas, “Scaling Violations in Inclusive e^+e^- Annihilation Spectra,” *Phys. Rev. D* **27** (1983) 105. doi:10.1103/PhysRevD.27.105
- [65] J. D. Bjorken, “Properties of Hadron Distributions in Reactions Containing Very Heavy Quarks,” *Phys. Rev. D* **17** (1978) 171. doi:10.1103/PhysRevD.17.171
- [66] D. Amati and G. Veneziano, “Preconfinement as a Property of Perturbative QCD,” *Phys. Lett. B* **83** (1979) 87. doi:10.1016/0370-2693(79)90896-7
- [67] D. J. Lange, “The EvtGen particle decay simulation package,” *Nucl. Instrum. Meth. A* **462** (2001) 152. doi:10.1016/S0168-9002(01)00089-4
- [68] S. Jadach, J. H. Kuhn and Z. Was, “TAUOLA: A Library of Monte Carlo programs to simulate decays of polarized tau leptons,” *Comput. Phys. Commun.* **64** (1990) 275. doi:10.1016/0010-4655(91)90038-M

- [69] S. Jadach, Z. Was, R. Decker and J. H. Kuhn, “The tau decay library TAUOLA: Version 2.4,” *Comput. Phys. Commun.* **76** (1993) 361. doi:10.1016/0010-4655(93)90061-G
- [70] J. Alwall, M. Herquet, F. Maltoni, O. Mattelaer and T. Stelzer, “MadGraph 5 : Going Beyond,” *JHEP* **1106** (2011) 128 doi:10.1007/JHEP06(2011)128 [arXiv:1106.0522 [hep-ph]].
- [71] S. Alioli, P. Nason, C. Oleari and E. Re, “A general framework for implementing NLO calculations in shower Monte Carlo programs: the POWHEG BOX,” *JHEP* **1006** (2010) 043 doi:10.1007/JHEP06(2010)043 [arXiv:1002.2581 [hep-ph]].
- [72] R. Frederix, S. Frixione, V. Hirschi, F. Maltoni, R. Pittau and P. Torrielli, “Four-lepton production at hadron colliders: aMC@NLO predictions with theoretical uncertainties,” *JHEP* **1202** (2012) 099 doi:10.1007/JHEP02(2012)099 [arXiv:1110.4738 [hep-ph]].
- [73] J. Alwall *et al.*, “The automated computation of tree-level and next-to-leading order differential cross sections, and their matching to parton shower simulations,” *JHEP* **1407** (2014) 079 doi:10.1007/JHEP07(2014)079 [arXiv:1405.0301 [hep-ph]].
- [74] P. Artoisenet, R. Frederix, O. Mattelaer and R. Rietkerk, “Automatic spin-entangled decays of heavy resonances in Monte Carlo simulations,” *JHEP* **1303** (2013) 015 doi:10.1007/JHEP03(2013)015 [arXiv:1212.3460 [hep-ph]].
- [75] M. L. Mangano, M. Moretti and R. Pittau, “Multijet matrix elements and shower evolution in hadronic collisions: $Wb\bar{b} + n$ jets as a case study,” *Nucl. Phys. B* **632** (2002) 343 doi:10.1016/S0550-3213(02)00249-3 [hep-ph/0108069].
- [76] M. L. Mangano, M. Moretti, F. Piccinini, R. Pittau and A. D. Polosa, “ALPGEN, a generator for hard multiparton processes in hadronic collisions,” *JHEP* **0307** (2003) 001 doi:10.1088/1126-6708/2003/07/001 [hep-ph/0206293].
- [77] R. Frederix and S. Frixione, “Merging meets matching in MC@NLO,” *JHEP* **1212** (2012) 061 doi:10.1007/JHEP12(2012)061 [arXiv:1209.6215 [hep-ph]].
- [78] N. D. Christensen and C. Duhr, “FeynRules - Feynman rules made easy,” *Comput. Phys. Commun.* **180** (2009) 1614 doi:10.1016/j.cpc.2009.02.018 [arXiv:0806.4194 [hep-ph]].
- [79] S. Frixione, Z. Kunszt and A. Signer, “Three jet cross-sections to next-to-leading order,” *Nucl. Phys. B* **467** (1996) 399 doi:10.1016/0550-3213(96)00110-1 [hep-ph/9512328].
- [80] S. Frixione, “A General approach to jet cross-sections in QCD,” *Nucl. Phys. B* **507** (1997) 295 doi:10.1016/S0550-3213(97)00574-9 [hep-ph/9706545].

- [81] V. Hirschi, R. Frederix, S. Frixione, M. V. Garzelli, F. Maltoni and R. Pittau, JHEP **1105** (2011) 044 doi:10.1007/JHEP05(2011)044 [arXiv:1103.0621 [hep-ph]].
- [82] S. Frixione and B. R. Webber, “Matching NLO QCD computations and parton shower simulations,” JHEP **0206** (2002) 029 doi:10.1088/1126-6708/2002/06/029 [hep-ph/0204244].
- [83] B. K. Bullock, K. Hagiwara and A. D. Martin, “Tau polarization and its correlations as a probe of new physics,” Nucl. Phys. B **395** (1993) 499. doi:10.1016/0550-3213(93)90045-Q
- [84] T. Sjostrand, S. Mrenna and P. Z. Skands, “A Brief Introduction to PYTHIA 8.1,” Comput. Phys. Commun. **178** (2008) 852 doi:10.1016/j.cpc.2008.01.036 [arXiv:0710.3820 [hep-ph]].
- [85] S. Agostinelli *et al.* [GEANT4 Collaboration], “GEANT4: A Simulation toolkit,” Nucl. Instrum. Meth. A **506** (2003) 250. doi:10.1016/S0168-9002(03)01368-8
- [86] S. Chatrchyan *et al.* [CMS Collaboration], “Description and performance of track and primary-vertex reconstruction with the CMS tracker,” JINST **9** (2014) 10, P10009 doi:10.1088/1748-0221/9/10/P10009 [arXiv:1405.6569 [physics.ins-det]].
- [87] S. Baffioni *et al.*, “Electron reconstruction in CMS,” Eur. Phys. J. C **49** (2007) 1099. doi:10.1140/epjc/s10052-006-0175-5
- [88] S. Chatrchyan *et al.* [CMS Collaboration], “Performance of CMS muon reconstruction in pp collision events at $\sqrt{s} = 7$ TeV,” JINST **7** (2012) P10002 doi:10.1088/1748-0221/7/10/P10002 [arXiv:1206.4071 [physics.ins-det]].
- [89] V. Khachatryan *et al.* [CMS Collaboration], JHEP **1101** (2011) 080 doi:10.1007/JHEP01(2011)080 [arXiv:1012.2466 [hep-ex]].
- [90] The CMS Collaboration: CMS Document 5581-v1
<https://cms-docdb.cern.ch/cgi-bin/PublicDocDB/ShowDocument?docid=5581>
 (29.11.2016)
- [91] [CMS Collaboration], “Particle-Flow Event Reconstruction in CMS and Performance for Jets, Taus, and MET,” CMS-PAS-PFT-09-001.
- [92] V. Khachatryan *et al.* [CMS Collaboration], “Jet energy scale and resolution in the CMS experiment in pp collisions at 8 TeV,” [arXiv:1607.03663 [hep-ex]].
- [93] S. Chatrchyan *et al.* [CMS Collaboration], “Missing transverse energy performance of the CMS detector,” JINST **6** (2011) P09001 doi:10.1088/1748-0221/6/09/P09001 [arXiv:1106.5048 [physics.ins-det]].

- [94] S. Chatrchyan *et al.* [CMS Collaboration], ‘Identification of b-quark jets with the CMS experiment,” JINST **8** (2013) P04013 doi:10.1088/1748-0221/8/04/P04013 [arXiv:1211.4462 [hep-ex]].
- [95] CMS Collaboration [CMS Collaboration], “Performance of b tagging at $\sqrt{s}=8$ TeV in multijet, ttbar and boosted topology events,” CMS-PAS-BTV-13-001.
- [96] CMS Collaboration [CMS Collaboration], “Identification of b quark jets at the CMS Experiment in the LHC Run 2,” CMS-PAS-BTV-15-001.
- [97] P. Billoir, “Progressive track recognition with a Kalman like fitting procedure,” Comput. Phys. Commun. **57** (1989) 390. doi:10.1016/0010-4655(89)90249-X
- [98] P. Billoir and S. Qian, “Simultaneous pattern recognition and track fitting by the Kalman filtering method,” Nucl. Instrum. Meth. A **294** (1990) 219. doi:10.1016/0168-9002(90)91835-Y
- [99] R. Mankel, “A Concurrent track evolution algorithm for pattern recognition in the HERA-B main tracking system,” Nucl. Instrum. Meth. A **395** (1997) 169. doi:10.1016/S0168-9002(97)00705-5
- [100] R. Frühwirth, “Application of Kalman filtering to track and vertex fitting,” Nucl. Instrum. Meth. A **262** (1987) 444. doi:10.1016/0168-9002(87)90887-4
- [101] R. Frühwirth, W. Waltenberger and P. Vanlaer, “Adaptive vertex fitting,” J. Phys. G **34** (2007) N343. doi:10.1088/0954-3899/34/12/N01
- [102] W. Adam, R. Frühwirth, A. Strandlie and T. Todor, “Reconstruction of Electrons with the Gaussian-Sum Filter in the CMS Tracker at the LHC”, J.Phys. G 31 (2005) N9, doi:10.1088/0954-3899/31/9/N01
- [103] S. Chatrchyan *et al.* [CMS Collaboration], “Energy Calibration and Resolution of the CMS Electromagnetic Calorimeter in pp Collisions at $\sqrt{s} = 7$ TeV,” JINST **8** (2013) P09009 [JINST **8** (2013) 9009] doi:10.1088/1748-0221/8/09/P09009 [arXiv:1306.2016 [hep-ex]].
- [104] M. Anfreville *et al.*, “Laser monitoring system for the CMS lead tungstate crystal calorimeter,” Nucl. Instrum. Meth. A **594** (2008) 292. doi:10.1016/j.nima.2008.01.104
- [105] A. Hocker *et al.*, “TMVA - Toolkit for Multivariate Data Analysis,” PoS ACAT (2007) 040 [physics/0703039 [PHYSICS]].
- [106] V. Khachatryan *et al.* [CMS Collaboration], Phys. Lett. B **766** (2017) 268 doi:10.1016/j.physletb.2017.01.011 [arXiv:1607.06943 [hep-ex]].

- [107] [CMS Collaboration], “Particle-flow commissioning with muons and electrons from J/Psi and W events at 7 TeV,” CMS-PAS-PFT-10-003.
- [108] S. Chatrchyan *et al.* [CMS Collaboration], “Measurement of the Underlying Event Activity at the LHC with $\sqrt{s} = 7$ TeV and Comparison with $\sqrt{s} = 0.9$ TeV,” JHEP **1109** (2011) 109 doi:10.1007/JHEP09(2011)109 [arXiv:1107.0330 [hep-ex]].
- [109] M. Cacciari and G. P. Salam, “Pileup subtraction using jet areas,” Phys. Lett. B **659** (2008) 119 doi:10.1016/j.physletb.2007.09.077 [arXiv:0707.1378 [hep-ph]].
- [110] CMS Collaboration [CMS Collaboration], “Search for a neutral MSSM Higgs boson decaying into $\tau\tau$ at 13 TeV,” CMS-PAS-HIG-16-006.
- [111] [CMS Collaboration], “Measurement of Tracking Efficiency,” CMS-PAS-TRK-10-002.
- [112] M. Cacciari, G. P. Salam and G. Soyez, “The Anti-k(t) jet clustering algorithm,” JHEP **0804** (2008) 063 doi:10.1088/1126-6708/2008/04/063 [arXiv:0802.1189 [hep-ph]].
- [113] M. Cacciari, G. P. Salam and G. Soyez, “The Catchment Area of Jets,” JHEP **0804** (2008) 005 doi:10.1088/1126-6708/2008/04/005 [arXiv:0802.1188 [hep-ph]].
- [114] S. Chatrchyan *et al.* [CMS Collaboration], “Study of the underlying event at forward rapidity in pp collisions at $\sqrt{s} = 0.9, 2.76$, and 7 TeV,” JHEP **1304** (2013) 072 doi:10.1007/JHEP04(2013)072 [arXiv:1302.2394 [hep-ex]].
- [115] V. Khachatryan *et al.* [CMS Collaboration], “Event generator tunes obtained from underlying event and multiparton scattering measurements,” arXiv:1512.00815 [hep-ex].
- [116] R. Field, “Early LHC Underlying Event Data - Findings and Surprises,” arXiv:1010.3558 [hep-ph].
- [117] V. Khachatryan *et al.* [CMS Collaboration], “Measurement of $B\bar{B}$ Angular Correlations based on Secondary Vertex Reconstruction at $\sqrt{s} = 7$ TeV,” JHEP **1103** (2011) 136 doi:10.1007/JHEP03(2011)136 [arXiv:1102.3194 [hep-ex]].
- [118] V. Khachatryan *et al.* [CMS Collaboration], “Event generator tunes obtained from underlying event and multiparton scattering measurements,” Eur. Phys. J. C **76** (2016) no.3, 155 doi:10.1140/epjc/s10052-016-3988-x [arXiv:1512.00815 [hep-ex]].
- [119] CMS Collaboration [CMS Collaboration], “Search for Higgs Boson Production in Association with a Top-Quark Pair and Decaying to Bottom Quarks or Tau Leptons,” CMS-PAS-HIG-13-019.
- [120] J. Beringer *et al.* [Particle Data Group], “Review of Particle Physics (RPP),” Phys. Rev. D **86** (2012) 010001. doi:10.1103/PhysRevD.86.010001

- [121] CMS Twiki: “Samples for BTV Performance Measurements”
<https://twiki.cern.ch/twiki/bin/viewauth/CMS/BTagPerformanceSamples>
 (11.03.2016)
- [122] ZEUS Collaboration: “Measurement of charm fragmentation fractions in photoproduction at HERA”, arXiv:1306.4862v1 [hep-ex] 2013
- [123] CLEO Collaboration: “Charm meson spectra in e^+e^- annihilation at 10.5 GeV c.m.e.”, arXiv:hep-ex/0402040v2 (2004)
- [124] https://twiki.cern.ch/twiki/pub/CMS/BtagRecommendation53XReReco/SFb-pt.NOttbar_payload_EPS13.txt (May 2016)
- [125] CMS Collaboration [CMS Collaboration], “Search for Lepton Flavour Violating Decays of the Higgs Boson in the mu-tau final state at 13 TeV,” CMS-PAS-HIG-16-005.
- [126] S. Chatrchyan *et al.* [CMS Collaboration], “Evidence for the 125 GeV Higgs boson decaying to a pair of τ leptons,” JHEP **1405** (2014) 104 doi:10.1007/JHEP05(2014)104 [arXiv:1401.5041 [hep-ex]].
- [127] CMS Collaboration [CMS Collaboration], “Measurements of inclusive and differential Z boson production cross sections in pp collisions at $\sqrt{s} = 13$ TeV,” CMS-PAS-SMP-15-011.
- [128] M. Czakon and A. Mitov, “Top++: A Program for the Calculation of the Top-Pair Cross-Section at Hadron Colliders,” Comput. Phys. Commun. **185** (2014) 2930 doi:10.1016/j.cpc.2014.06.021 [arXiv:1112.5675 [hep-ph]].
- [129] CMS Collaboration [CMS Collaboration], “Measurement of the top quark pair production cross section in proton-proton collisions at $\sqrt{s}=13$ TeV with the CMS detector,” CMS-PAS-TOP-15-003.
- [130] T. Gehrmann, M. Grazzini, S. Kallweit, P. Maierhöfer, A. von Manteuffel, S. Pozzorini, D. Rathlev and L. Tancredi, “ W^+W^- Production at Hadron Colliders in Next to Next to Leading Order QCD,” Phys. Rev. Lett. **113** (2014) no.21, 212001 doi:10.1103/PhysRevLett.113.212001 [arXiv:1408.5243 [hep-ph]].
- [131] CMS Collaboration [CMS Collaboration], “Measurement of the WW cross section pp collisions at $\sqrt{s}=13$ TeV,” CMS-PAS-SMP-16-006.
- [132] J. M. Campbell and R. K. Ellis, “MCFM for the Tevatron and the LHC,” Nucl. Phys. Proc. Suppl. **205-206** (2010) 10 doi:10.1016/j.nuclphysbps.2010.08.011 [arXiv:1007.3492 [hep-ph]].

- [133] CMS Collaboration [CMS Collaboration], “Measurement of the WZ production cross section in pp collisions at $\sqrt{s} = 13$ TeV,” CMS-PAS-SMP-16-002.
- [134] CMS Collaboration [CMS Collaboration], “Measurement of the ZZ production cross section in $\ell\ell'\ell'$ decays in pp collisions at $\sqrt{s} = 13$ TeV,” CMS-PAS-SMP-15-005.
- [135] G. Cowan, K. Cranmer, E. Gross and O. Vitells, “Asymptotic formulae for likelihood-based tests of new physics,” Eur. Phys. J. C **71** (2011) 1554 Erratum: [Eur. Phys. J. C **73** (2013) 2501] doi:10.1140/epjc/s10052-011-1554-0, 10.1140/epjc/s10052-013-2501-z [arXiv:1007.1727 [physics.data-an]].
- [136] [ATLAS Collaboration], “Procedure for the LHC Higgs boson search combination in summer 2011,” ATL-PHYS-PUB-2011-011, ATL-COM-PHYS-2011-818, CMS-NOTE-2011-005.
- [137] T. Junk, “Confidence level computation for combining searches with small statistics,” Nucl. Instrum. Meth. A **434** (1999) 435 doi:10.1016/S0168-9002(99)00498-2 [hep-ex/9902006].
- [138] A. L. Read, “Presentation of search results: The CL(s) technique,” J. Phys. G **28** (2002) 2693. doi:10.1088/0954-3899/28/10/313
- [139] CMS Collaboration [CMS Collaboration], “CMS Luminosity Measurement for the 2015 Data Taking Period,” CMS-PAS-LUM-15-001.
- [140] S. Chatrchyan *et al.* [CMS Collaboration], “Observation of the associated production of a single top quark and a W boson in pp collisions at $\sqrt{s} = 8$ TeV,” Phys. Rev. Lett. **112** (2014) no.23, 231802 doi:10.1103/PhysRevLett.112.231802 [arXiv:1401.2942 [hep-ex]].
- [141] S. Chatrchyan *et al.* [CMS Collaboration], “Measurement of $W\gamma$ and $Z\gamma$ production in pp collisions at $\sqrt{s} = 7$ TeV,” Phys. Lett. B **701** (2011) 535 doi:10.1016/j.physletb.2011.06.034 [arXiv:1105.2758 [hep-ex]].
- [142] <https://twiki.cern.ch/twiki/bin/view/LHCPhysics/CERNYellowReportPageAt13TeV> (06.12.2016)
- [143] V. Khachatryan *et al.* [CMS Collaboration], “Measurement of the differential cross section for top quark pair production in pp collisions at $\sqrt{s} = 8$ TeV,” Eur. Phys. J. C **75** (2015) no.11, 542 doi:10.1140/epjc/s10052-015-3709-x [arXiv:1505.04480 [hep-ex]].
- [144] A. Denner, S. Heinemeyer, I. Puljak, D. Rebuszi and M. Spira, “Standard Model Higgs-Boson Branching Ratios with Uncertainties,” Eur. Phys. J. C **71** (2011) 1753 doi:10.1140/epjc/s10052-011-1753-8 [arXiv:1107.5909 [hep-ph]].

- [145] <https://twiki.cern.ch/twiki/bin/view/LHCPhysics/CERNYellowReportPageAt8TeV> (19.01.2017)
- [146] V. Khachatryan *et al.* [CMS Collaboration], “Measurement of the t-tbar production cross section in the e-mu channel in proton-proton collisions at $\sqrt{s} = 7$ and 8 TeV,” JHEP **1608** (2016) 029 doi:10.1007/JHEP08(2016)029 [arXiv:1603.02303 [hep-ex]].
- [147] G. Aad *et al.* [ATLAS Collaboration], “Search for lepton-flavour-violating $H \rightarrow \mu\tau$ decays of the Higgs boson with the ATLAS detector,” JHEP **1511** (2015) 211 doi:10.1007/JHEP11(2015)211 [arXiv:1508.03372 [hep-ex]].
- [148] G. Aad *et al.* [ATLAS Collaboration], “Search for lepton-flavour-violating decays of the Higgs and Z bosons with the ATLAS detector,” arXiv:1604.07730 [hep-ex].
- [149] A. Elagin, P. Murat, A. Pranko and A. Safonov, “A New Mass Reconstruction Technique for Resonances Decaying to di-tau,” Nucl. Instrum. Meth. A **654** (2011) 481 doi:10.1016/j.nima.2011.07.009 [arXiv:1012.4686 [hep-ex]].
- [150] S. Bressler, A. Dery and A. Efrati, “Asymmetric lepton-flavor violating Higgs boson decays,” Phys. Rev. D **90** (2014) no.1, 015025 doi:10.1103/PhysRevD.90.015025 [arXiv:1405.4545 [hep-ph]].

Acknowledgements

I would like to thank all people who supported me during my time at the Universität Hamburg. First, I want to thank Peter Schleper for supervising my thesis and his support. Furthermore, I would like to thank him and Elisabetta Gallo for reading and evaluating my dissertation. Many thanks to the other members of my examination commission, Caren Hagner, Gudrid Moortgat-Pick, and Alexei Raspereza.

Also I would like to especially thank Adrian Perieanu, Thomas Peiffer, Daniel Tröndle, and Benedikt Vormwald for their advice and support. Special thanks to Adrian Perieanu, Thomas Peiffer, Benedikt Vormwald, Mikko Meyer, Laura Vanhoefer, and Gerrit Winkler for proofreading my entire thesis or parts of it.

In addition, I would like to thank my other colleagues and friends. Especially, I would like to thank Susanne Ott, Sebastian Laurien, Niklas Pietsch, Tobias Lapsien, Peter Drechsel, and Denis Rathjens.

This work would have not been possible without the support of my family. I want to thank my parents, Sabine and Jörg Vanhoefer, for always supporting me and for your advice. I also want to thank my brother and sisters, Tim, Lisa, and my twin Laura, thank you for always being there. I would also like to thank my friends Sabrina, Jeanny, Eva, Lena, and Fred for their support.

Finally, I would like to thank Gerrit for supporting me, especially in the last period before finishing my PhD.

List of Publications

Papers

CMS Collaboration: “Search for lepton flavour violating decays of the Higgs boson to $e\tau$ and $e\mu$ in proton-proton collisions at $\sqrt{s} = 8$ TeV,” Phys. Lett. B **763** (2016) 472 doi:10.1016/j.physletb.2016.09.062 [arXiv:1607.03561 [hep-ex]].

Conference Reports

Annika Vanhoefer: “Search for electron tau Lepton Flavour Violating Higgs Decay with CMS,” CMS-CR-2015-284 (November 2015)

Theses

- Master thesis: “Search for Supersymmetry at the LHC with Multivariate Analysis Techniques,” (March 2013)
- Bachelor thesis: “Untersuchung zur Datenqualität im OPERA-Experiment,” (March 2011)

Eidesstattliche Versicherung

Hiermit erkläre ich an Eides statt, dass ich die vorliegende Dissertationsschrift selbst verfasst und keine anderen als die angegebenen Quellen und Hilfsmittel benutzt habe.

Declaration on oath

I hereby declare, on oath, that I have written the present dissertation by my own and have not used other than the acknowledged resources and aids.

Hamburg, den 12. Mai 2017