

Search for new heavy charged gauge bosons with the ATLAS detector

by

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Ich versichere, dass ich die Arbeit selbstständig verfasst und keine anderen als die angegebenen Quellen und Hilfsmittel benutzt, sowie Zitate kenntlich gemacht habe.

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1. Introduction

Particle physics deals with the most fundamental level at which nature can be understood. Experiments in high energy conditions allow for precise study and measurements of the structure and dynamics of matter. In particle colliders, such as the Large Hadron Collider (LHC) at CERN unprecedented center-of-mass energies can be achieved in proton-proton collisions. By recording the products of these collisions, the processes that occurred can be reconstructed. The ATLAS detector allows for detection and storage of vast numbers of collision events, and offers precise measurements of properties of the particles produced in these events.

Using this approach, the Standard Model (SM) can be tested. This theoretical framework consists of a mathematical description of particle physics and allows for precise calculations and predictions. Even though the SM has proven to be highly successful, it does not explain everything observed in nature. Extensions of the SM can remedy this incompleteness. Using the same colliders, searches for these new phenomena can be conducted.

This thesis describes the search for new heavy charged gauge bosons. These bosons, collectively referred to as W' bosons, would appear as resonance structures on top of the expectation from the SM alone. Searches of this kind have been carried out in previous experiments, but 2015 saw the first collisions at the highest center-of-mass energy of 13 TeV achieved at the time of writing. This analysis uses the entire data collected in 2015 and 2016. Interference effects between the hypothesized W' bosons and the charged W gauge bosons of the weak interaction in the SM are considered in this analysis.

In the absence of a clear signal, upper limits on the relative coupling strength of new W' bosons can be calculated using a Bayesian statistics approach. These upper limits can be translated into a lower limit on the pole mass of the hypothesized boson under the assumption of a specific model. Below this pole mass the new particle can be excluded by the observed data at 95 % confidence level.

2. Theoretical background

2.1. The Standard Model of particle physics

The Standard Model (SM) of particle physics encompasses the set of theories, that is used to describe the world at the particle level. The SM describes three of the four forces in nature, the electromagnetic, the weak and the strong interaction. Each of these forces is responsible for different phenomena that can be observed in nature.

Matter consists of point like particles with an absolute spin of $1/2$, called *fermions*. There are two types of fermions, each subdivided into three *generations*. The first type are the *leptons*, consisting of the *electron*, *muon* and the *tau*, as well as a *neutrino* in each generation. The other type are the *quarks*. The three generations contain two quarks each, *up* and *down*, *charm* and *strange* and *top* and *bottom*. All of these fermions, except the neutrinos, carry an electromagnetic charge, and can interact via the electromagnetic force, which is mediated by the exchange of *photons*. Every fermion has an associated *antifermion*, with a flipped sign of the electromagnetic charge, and in the case of quarks, a flipped sign in the effective Baryon number of $B_q = 1/3$. The leptons (antileptons) carry integer electromagnetic charges of -1 ($+1$), while the quarks (antiquarks) carry fractional charges of $+2/3$ ($-1/3$).

Quarks (which will subsequently also refer to their antiparticles) can be bound to form *hadrons*, which are either pairs (mesons) or triplets (baryons). These hadrons are held together by the strong force, which is mediated by the exchange of *gluons*. The strong force couples to a color charge of red, green or blue. Quarks, antiquarks and gluons carry this charge, while hadrons are neutral. This also leads to an effect called *confinement*. If color neutral quark systems are separated, the strong binding increases. At some distance, the creation of additional quark-antiquark pairs is energetically preferential. The resulting system is then color neutral.

Heavy leptons, like the muon and the tau, and quarks can decay via the weak interaction, thereby producing a neutrino of the same generation, as well as a W boson, the exchange particle of the charged current weak interaction. Weak interaction can also occur via neutral current reactions, which are mediated by the Z boson. These interaction particles are called *bosons* because they carry integer spins of 1. The gluon and the photon are massless, while $W^\pm$ and Z are massive. This is explained by the *Higgs mechanism*, which also results in an additional spin-0 boson called the *Higgs*. The charges, spins and masses of the various particles are listed in Table 2.1.

2. Theoretical background

The SM provides a quantitative mathematical description of the properties of the various particles, as well as the interactions. The particles are described as *spinor* fields, encoding their spin as well as other properties. A common choice for a unit system in high energy particle physics is the natural one. The natural unit system sets the speed of light c and the reduced Planck constant $\hbar$ equal to 1, and is used throughout this thesis.

Table 2.1.: Properties of the fermions and bosons. The values are taken from the PDG [1]

name	charge	spin	mass [MeV]
quarks			
up	+2/3	1/2	2.2
down	-1/3	1/2	4.7
charm	+2/3	1/2	1.27×10^3
strange	-1/3	1/2	96
top	+2/3	1/2	173.21×10^3
bottom	-1/3	1/2	4.18×10^3
leptons			
electron	-1	1/2	0.51
electron neutrino	0	1/2	$< 2.2 \times 10^{-3}$
muon	-1	1/2	105.66
muon neutrino	0	1/2	< 1.7
tau	-1	1/2	1776.86
tau neutrino	0	1/2	< 15.5
bosons			
gluon	0	1	0
photon	0	1	0
Z	0	1	91.19×10^3
W	± 1	1	80.39×10^3
H	0	0	125.09×10^3

2.1.1. Mathematical structure

The mathematical structure which gives rise to the description of the particles found in nature is that of quantum field theories (QFT). The basis is a quantized description of fields $\phi_i(x_\mu)$ and their dynamics, which are encoded in a *Lagrangian* density $\mathcal{L}$. Particles are then identified as excited states of these underlying fields, and interactions are caused by coupling between these fields. The following introduction is based on references [2, 3].

The Euler-Lagrange formalism, known from classical mechanics, holds for Lagrangian densities. Hence, the Euler-Lagrange equation

$$\partial_\mu \left(\frac{\partial \mathcal{L}}{\partial(\partial_\mu \phi_i)} \right) - \frac{\partial \mathcal{L}}{\partial \phi_i} = 0 \quad (2.1)$$

yields field equations which describe the evolution of the corresponding fields. This applies to spin-0, spin-1 and spin-1/2 fields.

Spin-1

For a spin-1 field A_μ , the Lagrangian is

$$\mathcal{L} = -\frac{1}{16\pi} F^{\mu\nu} F_{\mu\nu} + \frac{1}{8\pi} m^2 A^\nu A_\nu \quad (2.2)$$

where

$$F^{\mu\nu} = \partial^\mu A^\nu - \partial^\nu A^\mu \quad (2.3)$$

is the field tensor. Making use of the Euler-Lagrange equation, this leads to

$$\partial_\mu F^{\mu\nu} + m^2 A^\nu = 0. \quad (2.4)$$

Spin-1/2

Starting from the Dirac-Lagrangian

$$\mathcal{L}_{\text{dir}} = i\bar{\Psi}\gamma^\mu\partial_\mu\Psi - m\bar{\Psi}\Psi \quad (2.5)$$

for a spin-1/2 field the dynamics for this type of particle can be derived. Ψ is the Dirac spinor of the particle, $\bar{\Psi} = \Psi^\dagger\gamma^0$ and γ^μ are the Dirac matrices. Using Equation (2.1) the Dirac equation and its adjoint

$$i\gamma^\mu\partial_\mu\Psi - m\Psi = 0 \quad i\partial_\mu\bar{\Psi}\gamma^\mu + m\bar{\Psi} = 0 \quad (2.6)$$

can be found for Ψ and $\bar{\Psi}$, respectively. Ψ and $\bar{\Psi}$ describe all fermions, and their antiparticles.

Local gauge invariance

In QFT, the concept that handles interactions is that of *local gauge symmetry*. The Dirac-Lagrangian $\mathcal{L}_{\text{dir}}$ fulfills *global gauge invariance* with respect to the symmetry group $U(1)$ without modification. This means that it remains invariant under

$$\Psi \rightarrow e^{i\theta}\Psi \quad (2.7)$$

where θ is an arbitrary phase angle. If, however, θ is a function of the position x^μ , $\mathcal{L}_{\text{dir}}$ is not invariant under

$$\Psi \rightarrow e^{i\theta(x)}\Psi = e^{-iq\lambda(x)}\Psi. \quad (2.8)$$

q is the charge of the Dirac particle. This breaking of the gauge symmetry $U(1)$ can be counteracted by introducing a coupling to a vector field A_μ . The Lagrangian then reads

$$\mathcal{L} = i\bar{\Psi}\gamma^\mu\partial_\mu\Psi - m\bar{\Psi}\Psi - (q\bar{\Psi}\gamma^\mu\Psi)A_\mu \quad (2.9)$$

2. Theoretical background

and is invariant under simultaneous local gauge transformations of Ψ and

$$A_\mu \rightarrow A_\mu + \partial_\mu \lambda(x). \quad (2.10)$$

The invariance is established by a redefinition of the derivative of the field Ψ to the *covariant derivative*

$$\mathcal{D}_\mu = \partial_\mu + iqA_\mu. \quad (2.11)$$

The Lagrangian from Equation (2.9) must also be supplemented with the free field components for A_μ , which can be found in Equation (2.2). This yields

$$\mathcal{L} = i\bar{\Psi}\gamma^\mu\partial_\mu\Psi - m\bar{\Psi}\Psi - (q\bar{\Psi}\gamma^\mu\Psi)A_\mu - \frac{1}{16\pi}F^{\mu\nu}F_{\mu\nu} + \frac{1}{8\pi}m_A A^\nu A_\nu. \quad (2.12)$$

Since the last term does not fulfill $U(1)$ invariance, the mass of the gauge field must be $m_A = 0$, which is compatible with observations. Excluding this term, Equation (2.12) becomes the Lagrangian of quantum electrodynamics (QED)

$$\mathcal{L}_{\text{QED}} = \bar{\Psi}(i\gamma^\mu\mathcal{D}_\mu - m)\Psi - \frac{1}{16\pi}F^{\mu\nu}F_{\mu\nu}, \quad (2.13)$$

where A_μ corresponds to the photon field. $\mathcal{L}_{\text{QED}}$ describes all electromagnetic phenomena.

In close analogy to the derivation of $\mathcal{L}_{\text{QED}}$, the Lagrangian describing the properties of the strong interaction can be found. As opposed to the QED case, the Lagrangian for quantum chromodynamics (QCD) is required to be invariant under transformations of the $SU(3)$ symmetry group. As this group has eight degrees of freedom, the local gauge invariance necessitates eight vector fields, corresponding to the gluons of different color charges. The Lagrangian for the strong interaction can be written as

$$\mathcal{L}_{\text{QCD}} = \bar{\Psi}(i\gamma^\mu\mathcal{D}_\mu^s - m)\Psi - \frac{1}{16\pi}\mathbf{F}_s^{\mu\nu}\mathbf{F}_{s\mu\nu}^s \quad (2.14)$$

with the strong covariant derivative $\mathcal{D}_\mu^s$ and the eight strong field tensors $\mathbf{F}_s^{\mu\nu}$, corresponding to the eight degrees of freedom of the gauge symmetry.

Spontaneous symmetry breaking and the electroweak theory

A problem arises in the formulation of an equivalent expression for the weak interaction. While the gauge fields from QED and QCD are massless, the gauge fields required in the weak interaction are not. With these heavy vector fields, the underlying symmetry is broken, and local gauge invariance is not achieved when considering the full Lagrangian. This fundamental issue can be addressed with spontaneous breaking of the underlying symmetry at low energies and the *Higgs* mechanism [4, 3]. The starting point is the understanding, that for the Lagrangians $\mathcal{L}_{\text{QED}}$ and $\mathcal{L}_{\text{QCD}}$, the ground state was found for vanishing fields $\phi = 0$. The ground state corresponds to the configuration of lowest energy. Assuming that the ground state does not vanish, but is offset by a certain margin through a potential $\mathcal{U}(\phi)$, the symmetry of the Lagrangian can be broken. Spontaneous symmetry breaking refers to the idea, that there is a local metastable minimum

$$\mathcal{U}(\phi = 0) = 0, \quad (2.15)$$

for which the Lagrangian is invariant. This potential can be extended to depend on a complex scalar field

$$\phi = v + \frac{1}{\sqrt{2}}(\phi_1 + i\phi_2) \quad (2.16)$$

with the potential

$$\mathcal{U}(\phi) = \mu^2(\phi^* \phi) + \lambda(\phi^* \phi)^2 \quad (2.17)$$

and the parameters μ and λ . This potential is then subtracted from the Lagrangian. For parameters $\lambda > 0$ and $\mu^2 < 0$, the ground state is now continuous around $\phi = 0$, with a vacuum expectation value of $\langle \phi \rangle = v \neq 0$. Nature effectively chooses an arbitrary configuration of ϕ_1 and ϕ_2 , which results in a specific ground state. In this ground state, the symmetry of the Lagrangian is then broken, and mass terms for the gauge fields are introduced.

Using these ingredients, the *Glashow-Weinberg-Salam* [5, 6, 7] model constructs a unified theory of the electromagnetic and weak interactions. Starting from a Lagrangian of massless fermions and requiring gauge invariance with respect to a $SU(2) \otimes U(1)$ symmetry, three gauge fields W_μ^i (corresponding to $SU(2)$) and an additional gauge field X_μ (corresponding to the $U(1)$ group of QED) are found. Further, a pair of complex scalar fields

$$\phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} \phi_3 + i\phi_4 \\ \phi_1 + i\phi_2 \end{pmatrix} \quad (2.18)$$

and a compatible potential $\mathcal{U}(\phi)$ are introduced. The Lagrangian, excluding the coupling of ϕ to the fermions, reads

$$\mathcal{L}_1 = (\mathcal{D}_\mu \phi)^* (\mathcal{D}_\mu \phi) - \frac{\mu^2}{2} \phi^* \phi - \frac{\lambda}{4} (\phi^* \phi)^2. \quad (2.19)$$

A convenient choice of ϕ leads to

$$\phi(x) = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ \sqrt{2}\eta + H(x) \end{pmatrix}. \quad (2.20)$$

The remaining components of ϕ can be forced to vanish by making use of the local invariance. The covariant derivative contains the four gauge fields W_μ^i and X_μ with coupling constants g and g' , respectively.

$$\begin{aligned} (\mathcal{D}_\mu \phi)^* (\mathcal{D}_\mu \phi) &= \frac{1}{2} (\partial_\mu H)^2 + \frac{g^2 \eta^2}{4} [(W_\mu^1)^2 + (W_\mu^2)^2] \\ &\quad + \frac{\eta^2}{4} (gW_\mu^3 - g'X_\mu)^2 \\ &\quad + \text{cubic} + \text{quartic terms}. \end{aligned} \quad (2.21)$$

In this form, the Lagrangian contains mass terms for the W_μ^i and X_μ fields. To produce a massless photon field A_μ , the fields W_μ^3 and X_μ can be rotated by an angle θ_W . This angle is called the *Weinberg angle*.

$$\begin{pmatrix} Z_\mu \\ A_\mu \end{pmatrix} = \begin{pmatrix} \cos \theta_W & -\sin \theta_W \\ \sin \theta_W & \cos \theta_W \end{pmatrix} \begin{pmatrix} W_\mu^3 \\ X_\mu \end{pmatrix} \quad (2.22)$$

In this rotation, W_μ^1 , W_μ^2 , Z_μ gain mass terms corresponding to masses of

$$M_{W_1}^2 = M_{W_2}^2 = M_W^2 = \frac{g^2 \eta^2}{2} \quad \text{and} \quad M_Z^2 = \frac{M_W^2}{\cos^2 \theta_W}, \quad (2.23)$$

2. Theoretical background

while A_μ is massless. The requirement of local gauge invariance stated above also introduces couplings of these four fields to the fermions. The complex scalar field ϕ also couples to fermions. A full Lagrangian $\mathcal{L}_2$ needs to include these coupling terms and the free field terms associated with the various fields in addition to $\mathcal{L}_1$. Considering the couplings and free field terms, W_μ^1 , W_μ^2 and Z_μ can be identified as the massive gauge bosons of the weak interaction. The massless vector field A_μ can be identified as the photon. The remaining scalar field $H(x)$ manifests itself as the Higgs boson, which was independently discovered at the LHC in the CMS [8] and ATLAS [9] experiments in 2012.

A peculiarity of the weak phenomenology is the violation of parity. Parity is violated by the neutral and charged currents, the charged current is *maximally* parity violating. The full Lagrangian for the interaction is of the *Vector minus Axial vector* ($V - A$) type. This manifests itself in a vertex factor of

$$\frac{-g_W}{2\sqrt{2}}\gamma^\mu(1 - \gamma^5) \quad (2.24)$$

where $\gamma^5 = i\gamma^0\gamma^1\gamma^2\gamma^3$. By considering the impact of the matrices on the spinor fields, it follows that the coupling term with γ^μ is a vector coupling, while the term $\gamma^\mu\gamma^5$ represents an axial vector coupling. As a consequence, the charged current only couples to *left-handed* particles and *right-handed* antiparticles, or rather their left- and right-handed components.

Observables in scattering processes

In order to test properties of the SM, experiments can be carried out. For an experiment to be useful, observable quantities need to be defined, which can then be measured. In the case of the scattering of particles, the primary observables are the rate Γ of a given decay and the probability for a process to take place. By counting the number of observed decays with a specific signature over a given time interval, the rate can be determined. Comparing the number of observed events with the total number of scattering processes that occurred, the probability for a process can be determined. The name of this probability is the *cross-section* σ of the process. In case a particle detector surrounds the interaction point completely, the relation

$$\frac{dN}{dt} = \sigma\mathcal{L} \quad (2.25)$$

yields the event rate, given the *luminosity* $\mathcal{L}$. The luminosity is the number of particles passing by the interaction point per unit time and unit area. If only a fraction $d\Omega$ of the full solid angle is covered by the detector, the differential cross-section is

$$\frac{d^2N}{dt d\Omega} = \frac{d\sigma}{d\Omega} \cdot \mathcal{L}. \quad (2.26)$$

For a scattering process of incoming particles 1, 2 and outgoing particles 3, 4 . . . n the cross-section is

$$\sigma = \frac{S}{4\sqrt{(p_1 p_2)^2 - (m_1 m_2)^2}} \int |\mathcal{M}_{fi}|^2 (2\pi)^4 \delta(p_1 + p_2 - p_3 - \dots - p_n) \times \prod_{j=3}^n 2\pi \delta(p_j^2 - m_j^2) \theta(p_j^0) \frac{d^4 p_j}{(2\pi)^4} \quad (2.27)$$

with the momenta p_i and the masses m_i of the various particles. $|\mathcal{M}_{fi}|^2$ is the scattering amplitude. S is a statistical factor, taking care of double-counting of identical particles. $S = 1$ if there are no identical particles in the final state. The cross-section depends on the matrix element

$$\mathcal{M}_{fi} = \langle f, t \rightarrow \infty | i, t \rightarrow -\infty \rangle \quad (2.28)$$

which encodes the transition from the initial state i to the final state f . The calculation of these matrix elements needs to be done for every combination of initial and final states. Using the *Feynman calculus*, the calculation can be carried out. It describes a method based on perturbation theory. The matrix element can then be written as a series with a term for every possible transition scenario between the states and each such transition scenario can be expressed by a *Feynman diagram*. The terms in the series can be categorized by the order of the coupling constant. It turns out, that the number of diagrams contributing to each transition is infinite and the series diverges. Higher order diagrams feature internal loops, or additional boson exchange between participating particles. Some examples of higher order diagrams are shown in the lower row of Figure 2.1. Calculations which only include terms linear in the coupling constant are referred to as leading order (LO). Examples for LO diagrams are found in the upper row of Figure 2.1.

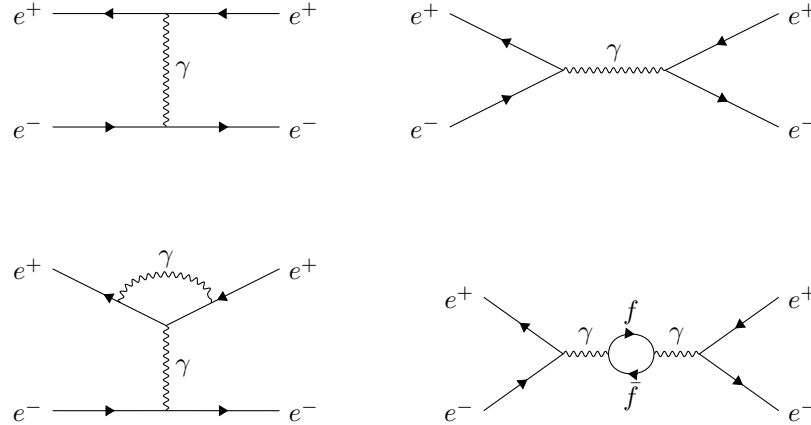


Figure 2.1.: Four Feynman diagrams as examples of configurations contributing to the matrix element for the process $e^+e^- \rightarrow e^+e^-$.

The divergence can be absorbed into a redefinition of the coupling constants and masses of the involved gauge fields using a sophisticated procedure called *renormalization*. The remaining series is still infinite, but each term has an increasing order of the relevant coupling constants g_i . As such, these higher order terms have diminishing impact on the total scattering amplitude, as long as $g_i \ll 1$. Nevertheless, computation of higher order terms is necessary for precise predictions of the scattering cross-sections. A consequence of higher orders is that after renormalization, the coupling constants become dependent on the energy scale of the process. The energy scale can be expressed as the momentum transfer Q^2 in a scattering process.

2. Theoretical background

This becomes problematic in the case of QCD. For high momentum transfers, $\alpha_s = g_s/4\pi$ is small, and the aforementioned perturbative calculation of the matrix element holds. However, α_s diverges at low momentum transfers Q^2 . If the coupling grows too large, the perturbative approach breaks down, since higher order terms are no longer suppressed by the coupling. Therefore, a precise perturbative calculation of the strong interaction is only possible at high energies, whereas at lower energies, phenomenological methods have to be employed.

This also explains the phenomenon of confinement. As the distance between quarks increases, the energy scale decreases. Consequently, the coupling constant increases, until the creation of additional quark-antiquark pairs yields an energetically lower state. Quarks in the final state of a scattering process will therefore not be detected as isolated particles, but as showers of hadrons that are the result of this hadronization process. This signature is commonly referred to as a *jet*.

2.1.2. Proton-Proton interactions and the parton model

In contrast to the fundamental particles, whose interactions were discussed in the previous section, protons are composite particles. The proton consists of three valence quarks, two *up* quarks and one *down* quark. These valence quarks determine the overall electromagnetic charge of +1 of the proton as well as the Baryon number of $B = +1$. Strong interaction via the exchange of gluons binds the proton system together. During gluon exchange, virtual additional gluons or quark-antiquark pairs can be produced in loops.

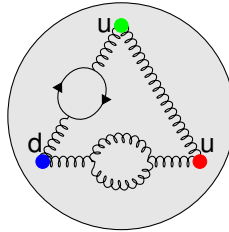


Figure 2.2.: Illustration of the composition of a proton and binding interactions.

In the *parton* model of the proton, the *valence* quarks, the virtual *sea* quarks, as well as the virtual gluons are referred to as partons. At any point in time, the total momentum of the proton is distributed among the partons. An illustration of this is shown in Figure 2.2. During a proton-proton collision with a hard scattering process, the assumption can be made that the fundamental interaction happens between a pair of partons from the incoming protons. In the case of high center-of-mass energies, the binding interactions between the partons can be neglected on the time scale of the scattering process. A sketch of the interaction can be seen in Figure 2.3. Since the precise internal configuration of the proton can not be calculated from QCD, the cross-section for a process including proton-proton interaction must be approximated. As described in [10], the cross section can be expressed by the cross-section of the interaction at the parton level, and factors which encode the structural information of the proton:

$$\sigma_{AB} = \int dx_a dx_b f_{a/A}(x_a, Q^2) f_{b/B}(x_b, Q^2) \hat{\sigma}_{ab \rightarrow X} \quad (2.29)$$

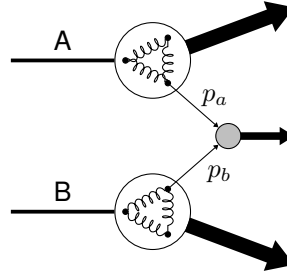


Figure 2.3.: Illustration of a proton-proton collision.

Here, x_a and x_b are the fractions of the total momentum of the proton which the partons a and b carry. The parton distribution functions (PDF) $f_{a/A}(x_a, Q^2)$ and $f_{b/B}(x_b, Q^2)$ describe the probability to find a parton a or b at a given momentum fraction x of the momentum of the proton A or B . The PDFs depend on the momentum transfer Q^2 of the scattering process, and are different for each type of parton. Finally, $\hat{\sigma}_{ab \rightarrow X}$ is the cross-section of the process at the parton level. By integrating over all possible momentum fractions, the cross-section at the proton level can be obtained. The separation of the cross-section of the reaction itself and the dynamics of the proton is described by the *factorization theorem* of QCD [11]. The PDFs can not be calculated from theory. To be able to get quantitative results for cross-sections involving protons, the PDFs have to be extracted from experimental data. Such measurements are performed by measuring the cross section of a known process at a fixed value of Q^2 . This can be performed by deep inelastic scattering in an electron-proton collider, fixed target experiments or from proton-proton collisions. From a cross-section measurement, the PDF for the proton can be extracted. Electron probing has the advantage that only one PDF is included in the cross-section, while proton-proton collisions can be performed at a wider range of Q^2 . A functional form for the cross-section including a parametrization of the PDF is fitted to the measured cross-section. This PDF can be extrapolated to other values of Q^2 by using the DGLAP¹ [12, 13, 14] equations.

An example of extracted PDFs for a range of x and $Q^2 = 10 \text{ GeV}^2$ and 10^4 GeV^2 is shown in Figure 2.4. PDFs for the various types of quarks and antiquarks, as well as for gluons, are shown. All PDFs decrease drastically towards higher values of x . This expresses the lower probability to find the proton in a state where a very large fraction of its momentum is carried by a single parton. For the up and down quarks there is a pronounced peak visible in the PDFs at $Q^2 = 10 \text{ GeV}^2$. These peaks are a consequence of up and down being the valence quarks of the proton. $\bar{u}$ and $\bar{d}$ are significantly less likely to be present for all values of x , since they have to be created in the *sea*. For charm and strange, there is no apparent difference between quarks and antiquarks, because both types are only found in the *sea*. Since the gluon PDF is significantly larger than the others, it is scaled by a factor of 1/10 for visibility. The peak structure of up and down quarks becomes less pronounced at $Q^2 = 10^4 \text{ GeV}^2$. At higher energies, resolving quarks and gluons from the *sea* becomes more likely. The plots also show uncertainties on the PDFs. These account for uncertainties introduced during the fitting procedure, as well as experimental uncertainties from the input data used in the extraction. There are dedicated groups studying the PDFs and PDF uncertainties, and how to propagate them into simulations of proton-proton collisions.

¹ Dokshitzer-Gribov-Lipatov-Altarelli-Parisi

2. Theoretical background

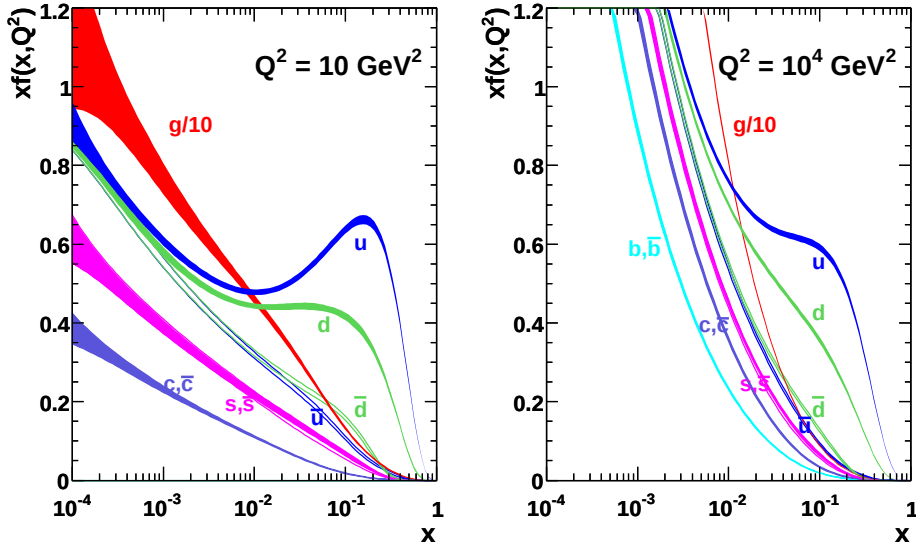


Figure 2.4.: Plots showing the parton distribution function for the different types of quarks and anti-quarks, as well as gluons. The gluon distribution has been scaled down for visibility. The dependence on x is shown for two values of Q^2 . Figures taken from [15].

2.2. Physics beyond the Standard Model

Even though the Standard Model, outlined in the previous sections, has proven to be highly successful, there are remaining questions.

As described before, the SM contains a consistent description of the electromagnetic and weak interactions in the form of the electroweak (EW) unification. The strong interaction, however, remains entirely separate. All attempts to unify the electroweak and the strong interactions in a Grand Unified Theory (GUT), for instance using spontaneous symmetry breaking, have turned out to be unsuccessful. Additionally, the fourth fundamental force in nature, gravity, is missing entirely in the SM. A consistent description of gravity in the framework of a quantum field theory has yet to be found. Lacking that, unification of the electroweak and strong sector with gravity is also an outstanding issue.

Another aspect has to do with neutrinos. While their weak interaction is described by the electroweak theory, the SM does not necessitate the inclusion of a Higgs-induced mass term. Neutrinos, however, have been shown to undergo oscillation. In various experiments, measuring solar neutrinos and neutrinos from nuclear reactors, a time dependent transition between the three neutrino flavours has been observed. This can be modeled by a mixing of these flavour eigenstates with a corresponding set of mass eigenstates by the PMNS² matrix [16, 17]. This oscillation, however, is only possible if the neutrinos do not have vanishing masses. A conclusive formalism which explains the masses of the neutrinos has yet to be found.

²Pontecorvo–Maki–Nakagawa–Sakata

Observations of the universe show a severe asymmetry between the amount of matter and antimatter. It has been ruled out that this *Baryon* asymmetry might be a local effect [18]. An explanation for the preference of matter over antimatter could be found in cosmology, but would most likely have implications at the fundamental particle level. The observed Baryon asymmetry would require several conditions to be true, one of which is the violation of the charge-parity (CP) symmetry. CP symmetry is violated in the SM, albeit not to the required degree necessary.

Through the observation of rotational velocity in galaxies, a discrepancy to the expectation from Newtonian mechanics is found. If there is only visible matter this discrepancy has to be resolved by modifying the theoretical description of gravity. Another possible solution is the inclusion of additional invisible matter. This *dark matter* cannot interact with ordinary matter via the electromagnetic force to a significant degree, thus avoiding observation. By introduction of additional dark matter, the rotational curves of stars in galaxies can be brought back to agreement with Newtonian gravity. Additional hints at dark matter are found in structure formation during the evolution of the universe, as well as in the cosmic microwave background (CMB). The latter is background radiation emitted early on in the history of the universe. It decoupled from matter when the temperature of the universe dropped below the energy threshold required for binding of atoms. At the time of emission, the baryonic matter density was oscillating, induced by gravitational interaction. The observation of anisotropies in the CMB gives insight into this oscillation, and is consistent with the existence of dark matter. Dark matter would have to be made up of particles, but all particles of the SM have been ruled out as candidates of dark matter. A model connecting dark matter to the SM would have to explain the very weak coupling and reproduce dark matter densities which are consistent with observations of the universe.

The light from distant stars, galaxies and supernovae can be observed. By comparing the wavelengths of the incoming light to the emission spectrum expected for the type of body, the *redshift* can be determined. The measurements show, that light from all objects is redshifted, meaning that they move away from the observer. This result is consistent with an expansion of the universe itself, where every point is seemingly at the very center of the expansion. Further study showed that the redshift increases the further away the observed object is located. Consequently, the expansion of the universe must be accelerating. A concept called *dark energy* can explain this expansion. The origin of this dark energy is yet to be understood.

Considering these four examples of phenomena outside of the explanation of the SM, it seems reasonable to search for extensions of the SM. These Beyond Standard Model (BSM) theories might be able to solve some or all of the outstanding issues, and also reproduce the very precise results of the SM itself. One approach is the search for new symmetry groups, and construction of Lagrangians which are invariant under them. Through careful examination of the properties of the symmetry groups and the interactions, the SM could be constructed after breaking of the initial symmetry. Theories like this would then contain additional symmetries aside from the ones in the SM. A way to test new physics is to look for signatures of the gauge fields that are a result of these additional symmetries.

2. Theoretical background

2.2.1. New heavy charged gauge bosons

To address the issues mentioned above, different approaches can be chosen. For the example of a GUT the starting point can be to find a symmetry group for which the unified Lagrangian is invariant. The theory has to recover the SM interactions, and typically contains additional symmetries. Due to the presence of these symmetries, some of the models predict heavy charged gauge bosons. Bosons like this are commonly referred to as W' in analogy to their SM counterparts.

Little Higgs models

As described in reference [19], one approach is to use an initial symmetry which breaks into multiple $SU(2) \otimes U(1)$ groups. By combining the resulting groups in a convenient way and gauging them correspondingly, these models can recover the unbroken electroweak symmetries. They also manage to retain the Higgs fields required for the electroweak symmetry breaking, while using a Higgs-type mechanism for the additional gauge fields. The symmetry breaking has to be conducted in a way that the corrections arising from the subgroups compensate each other to retain only a manageable divergence. This is referred to as *collective symmetry breaking*.

Models of this kind are consistent with a light Higgs boson, and new physics is hidden below the symmetry breaking scale at $\mathcal{O}(10\text{TeV})$. The *littlest Higgs* model is an example which uses $SU(5)$ as a starting symmetry. This symmetry can be separated into

$$SU(5) \supset SO(5) \otimes [SU(2) \otimes U(1)]_1 \otimes [SU(2) \otimes U(1)]_2 \quad (2.30)$$

which is then used for the elaborate symmetry breaking procedure required to separate the SM structures from the additional BSM effects. The model recovers the SM Higgs boson, and adds a new triplet of $W'^{\pm}$ and Z' gauge bosons. Another result is the addition of a new quark T .

Left-Right Symmetric Model

An unexplained property of the SM is the maximal parity violation of the charged current weak interactions. The Left-Right Symmetric Model (LRSM) adds an additional symmetry group to the ones needed for the known SM interactions:

$$SU(2)_R \otimes SU(2)_L \otimes U(1). \quad (2.31)$$

The $SU(2)_L \otimes U(1)$ symmetries form the observed weak interactions, which only couple to left-handed fermions. $SU(2)_R$ on the other hand leads to a new pair of charged bosons $W_R'^{\pm}$, which are right-handed and only couple to right-handed fermions. When considering the entire symmetry, parity is not violated. The model also explains the very small masses of the left-handed neutrinos, by requiring that the right-handed neutrinos be very heavy. The mass of the W_R is predicted to be significantly

2.3. Search strategy for new heavy charged gauge bosons

higher than the left-handed counterpart, explaining a lack of detection in collider experiments thus far. The impact of the left- or right-handedness of a new W' on the search strategy will be discussed in Section 2.3.1.

Exceptional group models

The *exceptional group* E_6 can be used as a starting point for a GUT as described in [19]. It features a number of decomposition paths, two of which are

$$\begin{aligned} E_6 \supset SO(10) &\rightarrow SU(4) \otimes SU(2)_L \otimes SU(2)_R \\ &\rightarrow SU(3) \otimes SU(2)_L \otimes SU(2)_R \otimes U(1). \end{aligned} \quad (2.32)$$

If one of these paths is chosen by nature, the subgroups $SU(2)_L \otimes SU(2)_R$ appear. Hence, the aforementioned LRSM can be contained in an E_6 model.

2.2.2. The Sequential Standard Model benchmark

To avoid testing for individual models containing W' bosons, a benchmark model can be used. In the case of this analysis, the Sequential Standard Model (SSM) [20] delivers the signal model that is searched for. This benchmark introduces three new gauge bosons, $W'^{\pm}$ and Z' , with identical couplings to the SM particles. For this analysis, the coupling of the W' to the SM W and Z is neglected. The masses of these new bosons do not mirror their SM counter parts, but are significantly higher. In the framework of the SSM, the W' can arbitrarily be set to be left- or right-handed.

2.3. Search strategy for new heavy charged gauge bosons

This search is carried out in the final state $l^+\nu_l$ or $l^-\bar{\nu}_l$, where l stands for electrons and muons. In proton-proton collisions, this final state can directly be produced by the processes illustrated in Figure 2.5.

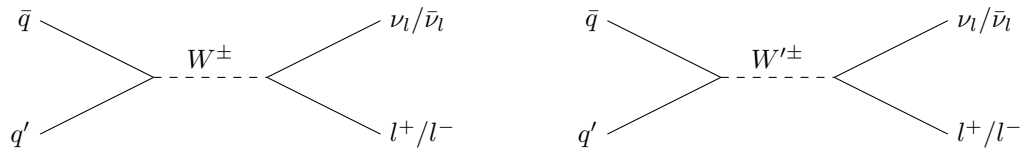


Figure 2.5.: Tree level diagrams showing the production of the $l\nu_l$ final state via a W' and a W boson.

A common search strategy for the discovery of a resonantly produced particle makes use of the invariant mass of its decay products. A search for resonance signatures introduced by the new

2. Theoretical background

particle can be conducted. A narrow resonance of a few percent relative width is assumed. As a consequence, the true width of the resonance is not expected to be resolved by the detector. When only including the narrow resonance of the signal process, the invariant mass distribution has a Breit-Wigner shape.

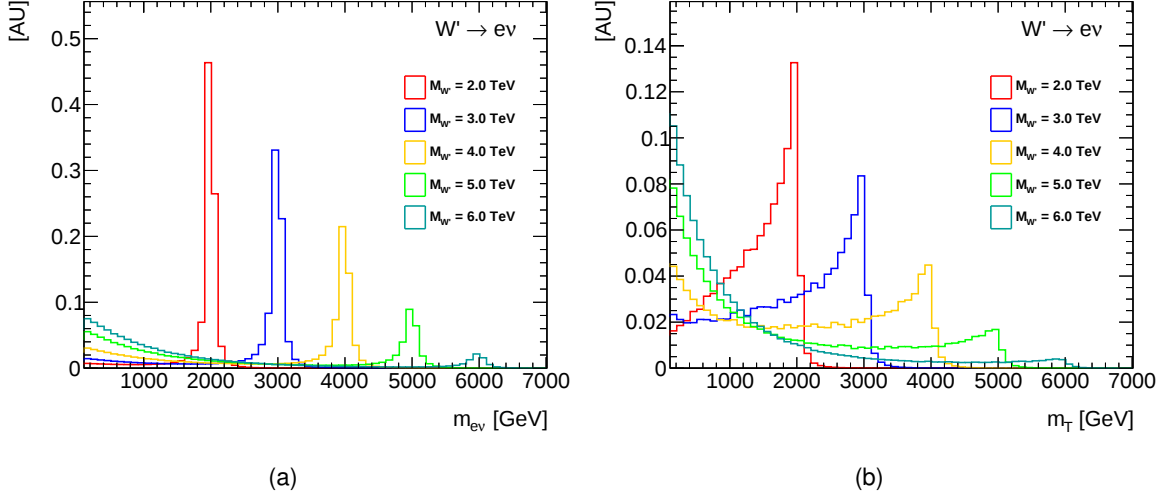


Figure 2.6.: Signal signatures in of the decay $W' \rightarrow e\nu_e$ for the invariant mass and the transverse mass for various pole masses of the W' . The signals are scaled to unit area to allow for a better comparison.

Invariant mass distributions for the signal process can be seen in Figure 2.6a for various pole masses. The resonance is relatively symmetric and centered around the pole mass value. A rise towards lower masses is observed, which becomes more pronounced for higher pole masses. This *parton-parton-luminosity* tail is a consequence of the proton-proton scattering process. The interacting partons only carry a momentum fraction whose probability density is described by the PDFs (see Section 2.1.2). The Breit-Wigner form of the cross-section is therefore convoluted with the PDF.

The leptonic final states of W and W' decays contain a neutrino. Due to their lack of charge and the fact that they only interact weakly, detecting neutrinos is a challenge. Dedicated neutrino detectors exist, which make use of significant amounts of material inside the detector volume in order to increase the probability of the neutrino to interact while passing through it. The products of the neutrino reaction can then be measured. The ATLAS detector is not specialized for measurement of neutrinos, hence the probability of detecting the actual neutrino from a W' event is negligible. While the neutrino's properties can not be measured in the detector, they still carry momentum. Conservation of momentum in the transverse plane of the detector can be used to recover some of the neutrino's properties. The calculation of this missing transverse momentum (E_T^{miss}) is described in Section 3.2.5. Using the transverse momentum of the charged lepton p_T^l and the transverse momentum of the neutrino p_T^ν , the transverse mass m_T can be calculated:

$$m_T = \sqrt{2p_T^l p_T^\nu [1 - \cos(\Delta\phi_{p_T^l, p_T^\nu})]} \quad (2.33)$$

Here, $\Delta\phi_{p_T^l, p_T^\nu}$ is the angle between the particle momenta in the transverse plane. The definition of m_T is analogous to the invariant mass, but only incorporates the transverse momentum components.

It is clear, that the transverse mass is only equal to the invariant mass, if the W' is produced at rest and if the decay takes place entirely in the transverse plane. Any other configuration results in a lower m_T . As seen in Figure 2.6b, the resonance is thus much broader and has a higher tail towards lower transverse masses. The tail above the pole mass stems from W' with transverse momentum due to initial state radiation and off-shell production of the W' .

As mentioned before, the neutrino momentum p_T^ν must be replaced by E_T^{miss} . A search based on this variable is therefore sensitive to any process which produces a lepton and missing transverse momentum. The dominant contribution is the off-shell production of the SM W . Background processes which have additional particles in the final states can be suppressed in the event selection.

2.3.1. Interference between SSM W' bosons and the Standard Model background

As discussed in Section 2.2.1, the helicity of additional W' particles varies by model. The LRSM predicts a right-handed particle, while other theories can accommodate both left- or right-handed scenarios. The helicity impacts the expected signature of the new boson. The helicities of the particles in the final state cannot be resolved by the detector.

If the proposed W' is itself right-handed, the coupling to fermions is exactly complementary to the SM W coupling. This means that the decays of W and W' into their final states $l\nu_l$, and all higher-order diagrams are actually different transitions, since they result in different helicity configurations. As a consequence, the amplitudes for the transitions do not mix. The lepton-neutrino pairs produced by the SM W are completely independent from the lepton-neutrino pairs produced by the W' . The diagrams shown in Figure 2.5 effectively contribute to separate matrix elements $[\bar{q}q' \rightarrow l\nu_l]_R$ and $[\bar{q}q' \rightarrow l\nu_l]_L$.

If, on the other hand, the W' is left-handed, the diagrams shown in Figure 2.5 are part of the exact same overall transition and are essentially indistinguishable. Calculation of the matrix element for $[\bar{q}q' \rightarrow l\nu_l]_L$ involves terms including both the SM W as well as the W' . This phenomenon is called interference between these two fields. Following Accomando et al. [21], the matrix element can be written as

$$|\mathcal{M}|^2 = |\mathcal{M}_{\text{SM}} + \mathcal{M}_{\text{SSM}}|^2 \quad (2.34)$$

$$= |\mathcal{M}_{\text{SM}}|^2 + |\mathcal{M}_{\text{SSM}}|^2 + 2 \text{Re}(\mathcal{M}_{\text{SM}}^* \cdot \mathcal{M}_{\text{SSM}}) \quad (2.35)$$

where $2 \text{Re}(\mathcal{M}_{\text{SM}}^* \cdot \mathcal{M}_{\text{SSM}})$ is the mixed-term that arises from the quadrature. In the first scenario, where W and W' couple exclusively, the mixed-term vanishes. Since the detection of the final state particles is entirely independent of the helicity states, the two processes are overlayed on top of each other. The number of events observed in this final state is thus only enhanced by the additional term $|\mathcal{M}_{\text{SSM}}|^2$. In the second scenario, the final state is identical, and the mixed-term does not vanish. Depending on the sign of the coupling constants, $2 \text{Re}(\mathcal{M}_{\text{SM}}^* \cdot \mathcal{M}_{\text{SSM}})$ might be purely positive. In

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this case, as in the previous scenario, the number of observed events is exclusively enhanced by the new particle. This is referred to as constructive interference. If the mixed-term is negative, on the other hand, the number of observed events in the final state might in fact be reduced, when compared to the expectation from the background only $|\mathcal{M}_{\text{SM}}|^2$ term. In this case, destructive interference is observed.

Both scenarios mentioned before will be considered. Using a reweighting approach, the mixed-term can be toggled on and off. In the case of a non-vanishing mixed-term, only destructive interference will be considered. Constructive interference would result in a higher number of signal events than without any interference. Therefore, the latter scenario predicts a smaller signal and is thus more conservative. A comparison between results with and without interference effects will be presented.

2.3.2. Relevant background processes

Signatures which are compatible with the signal final state discussed in Section 2.3 can come from a variety of processes. These background contributions can be reducible if their final state differs from the signal final state. Irreducible backgrounds are decays which produce a final state that is indistinguishable from the signal. Some processes are similar to the signal final state but produce additional jets. Since a large number of jets is expected in a hadron collider, these additional jets might not always be associated with their correct origin. Hence, such processes can also contribute background events.

Drell-Yan processes

In the context of a search for W' decaying to a lepton-neutrino pair, the decay of SM W poses the largest background. This is called the charged current Drell-Yan process [22]. Electrons, muons and taus are considered for this background even though the final state channels under study are restricted to electrons and muons only. Since it produces the exact same final state as the W' decay, except for helicity (see Section 2.3.1), the signature looks identical. The only difference between the two processes arises from the different pole masses. The SM W is produced resonantly at its pole mass of $M_W = 80.39 \text{ GeV}$. Off-shell production at higher masses is suppressed but still relevant. The SSM W' mass is set at various significantly higher values, meaning that the resonant production occurs in different mass ranges.

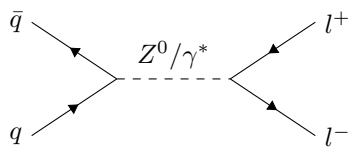


Figure 2.7.: Tree level diagram of the neutral current Drell-Yan process.

Another possibility for background is the production of a similar final state, where detector effects cause the signature to look like the final state of interest. The neutral current Drell-Yan process occurs by production of a virtual Z^0 or photon γ^* from quark-antiquark annihilation. The neutral boson can then decay into a lepton-antilepton pair. A Feynman diagram at tree level is shown in Figure 2.7. The final state of this process does not contain a neutrino. If, however, one of the produced leptons escapes the acceptance of the detector, the signature seen by the detector looks identical. The background can be suppressed by requiring events to have a transverse mass above the on-shell production of the W and Z bosons, and by vetoing events with additional leptons.

Diboson processes

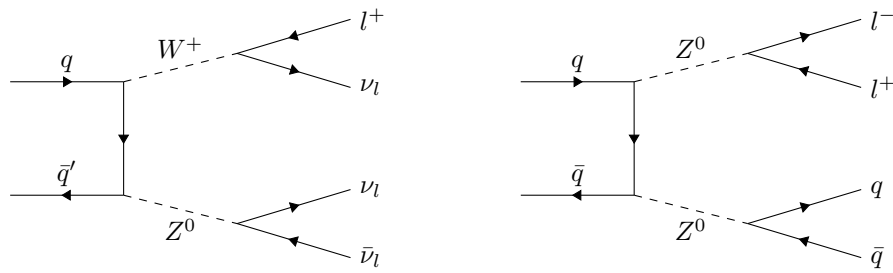


Figure 2.8.: Two example Feynman diagrams in the Diboson background process.

Two bosons can be produced at the same time by the exchange of a quark. The produced set of bosons is a combination of Z^0 and $W^\pm$, depending on the types of the incoming quarks. Each of the produced bosons decays individually, resulting in a variety of different final states. These final states can be fully leptonic, but also semileptonic. The final state $l\nu\nu\nu$ most closely resembles the signal final states, with only one lepton. In the case of $l\nu l\nu$ or $ll\nu\nu$, the lepton-lepton pair can produce a signal signature similar to the neutral current Drell-Yan process. If the decay is semileptonic, the quark signatures can easily be compatible with background jet activity. The remaining leptonic leg of the decay is then similar to the signal final state. Examples of Feynman diagrams of possible Diboson scenarios are shown in Figure 2.8.

Top processes

There are various processes which produce top quarks in their final states. In contrast to all other quarks, top quarks decay before hadronization. This means that nearly all produced top quarks decay into a bottom quark and a W (see Figure 2.9). The W can then either decay hadronically or leptonically. The latter case is a background source when analysing the lepton neutrino final state of the signal process.

Single top quarks can be produced (example diagrams in Figure 2.10), and processes producing a top in combination with a W (example diagrams in Figure 2.11) also contribute to the background.

2. Theoretical background

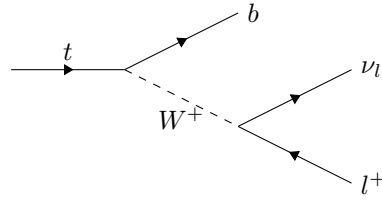


Figure 2.9.: Example Feynman diagram showing the decay of a top quark into a bottom quark and a W .

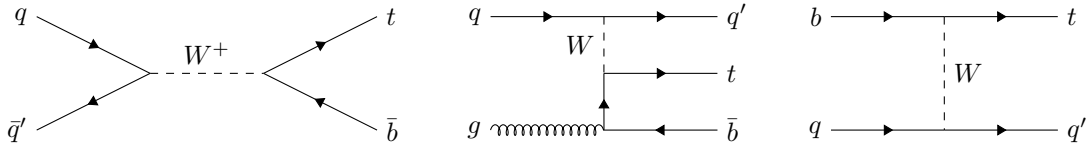


Figure 2.10.: Example Feynman diagrams showing the production of single top and one or more quarks.

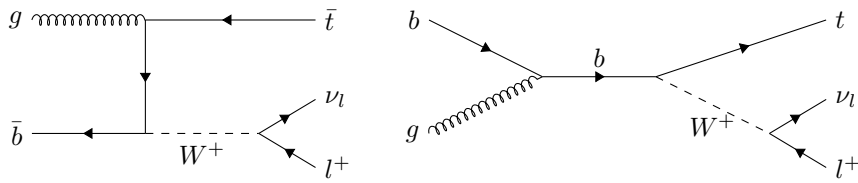


Figure 2.11.: Example Feynman diagrams showing the production of a top in association with a W boson.

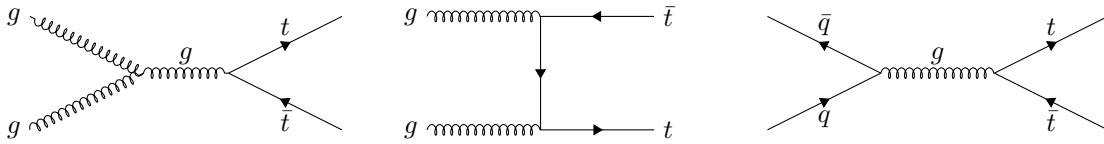


Figure 2.12.: Example Feynman diagrams showing the production of a pair of top quarks.

Finally, the pair production of top and antitop quarks can happen, for instance via gluon fusion. Examples for processes producing $t\bar{t}$ pairs are shown in Figure 2.12. Regardless of the production mechanism, the signatures of the final states of processes containing tops can be similar to the signal signature.



Figure 2.13.: Example Feynman diagrams showing the production of a pair of gluons.

Multijet background

The final background is closely related to the detection and reconstruction of particles in the experiment. There are numerous processes which result in a varying number of hadronic particles in the final state. An example for this is the production of gluon pairs. Since gluons couple to other gluons, the production of gluons in the final state is common. Two example diagrams of such processes are shown in Figure 2.13. The gluons produce jets that are reconstructed in the detector. These jets are susceptible to misidentification, and can produce signatures similar to the signal signature.

3. Experimental setup

In high energy particle physics, the goal of experiments is to produce conditions at which phenomena of interest take place. The main components of these experiments are particle accelerators and detectors. The former produce particle beams at high energies, in colliders these beams are made to collide with each other. Detectors can be set up around the collision point, in order to measure the products of the collisions.

3.1. The Large Hadron Collider

The Large Hadron Collider (LHC) is a proton-proton collider, and is at the time of writing part of the highest energy accelerator complex in the world. It is located at the European Organization for Nuclear Research (CERN) near Geneva, Switzerland. The collider itself is built inside a circular tunnel of about 27 km circumference below the ground.

The components of the accelerator complex are shown in Figure 3.1. It consists of a series of accelerators, most of which have been or are still used for other physics projects. The complex serves to bring bunches of protons to the target beam energy of 6.5 TeV and is designed to reach beam energies as high as 7 TeV. The proton bunches are initially brought to 50 MeV in the linear accelerator *LINAC2*¹. At the end of this accelerator, the protons are fed into the *BOOSTER*, *PS*² and *SPS*³, where they are accelerated to 1.4 GeV, 26 GeV and 450 GeV, respectively. After these three stages, the beams of proton bunches are extracted from their circular paths. The protons are fed into the two parallel beam pipes of the LHC itself. The beam pipes are enclosed by 1232 superconducting dipole magnets, which keep the protons on their circular trajectory around the center of the LHC. Additionally, 392 quadrupole focusing magnets are used to produce the required beam shapes. Since the protons travel in opposite directions, their trajectories have to be bent differently. To that end, the dipole magnets are designed so that the magnetic field lines in one beam pipe are antiparallel relative to the other beam pipe. In the LHC the protons are further accelerated up to their final energy. The machine was operated at 3.5 TeV from its initial startup in 2009 until 2011, when the beam energy was increased to 4 TeV. The data taking period up to and including 2012 is referred to as Run 1. The subsequent Run 2 data taking period started in 2015 and is scheduled to

¹Linear Accelerator 2

²Proton Synchrotron

³Super Proton Synchrotron

3. Experimental setup

continue until autumn of 2018. The beam energy was increased to 6.5 TeV for Run 2. At four points, the beams are bent by special dipole magnets so that they cross and interact with a certain *crossing angle*. Additionally, quadrupole magnets are used to focus the beams on the interaction point. The time gap between two subsequent crossings of bunches was 25 ns during the operation in 2015 and 2016. The available center-of-mass energy of collisions at the proton level is $\sqrt{s} = 13$ TeV in this configuration. This analysis uses data from 2015 and 2016.

Four major experiments have been constructed at the LHC, each of them is located at one of the four beam crossing points. The crossing positions are indicated in Figure 3.1. The CMS and ATLAS detectors are general purpose detectors, designed to cover a broad range of physics scenarios. LHCb is focused on studying interactions which result in events containing bottom quarks. In addition to the proton-proton collisions, the ALICE experiment studies collisions of lead nuclei, which are accelerated in the LHC in dedicated heavy ion periods. The analysis in this thesis uses data collected with the ATLAS experiment.

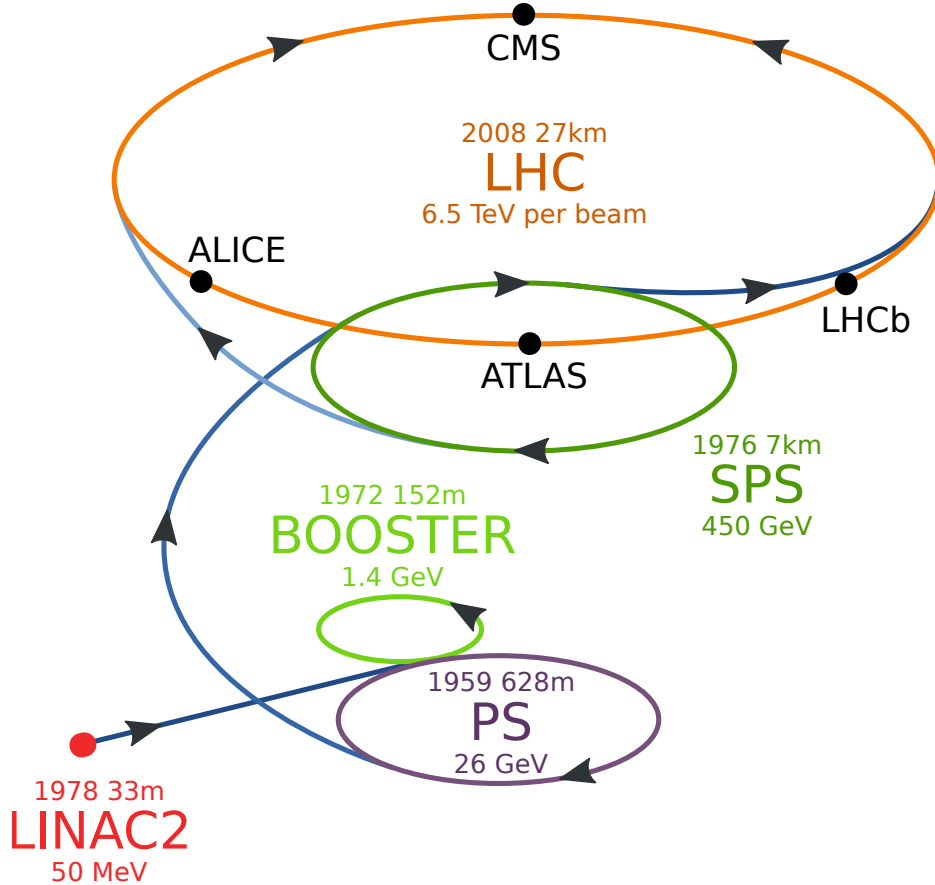


Figure 3.1.: The LHC accelerator complex at CERN. Figure slightly modified from [23].

A measure for the performance of a collider is the *luminosity*. Generally, the luminosity can be calculated as

$$\mathcal{L} = \frac{N^2 f N_b}{4\pi\sigma_x\sigma_y} \quad (3.1)$$

where f is the revolution frequency, N is the number of particles per bunch and N_b is the number

of bunches per beam. The expression assumes two beams with equal Gaussian profiles, which collide frontally without a crossing angle. $\sigma_{x,y}$ are the overlaps of the beams in x and y (see [24]). The interaction area $4\pi\sigma_x\sigma_y$ can also be expressed differently via the transverse emittance ϵ_n and β^* which characterize the spread and transverse extent of the beams. In combination, they effectively describe the brightness of the beam. Additionally, correction factors for the crossing angle and varying focussing at the crossing have to be included.

The design luminosity of the LHC is $10^{34} \text{ cm}^{-2} \text{ s}^{-1}$, and was first reached and surpassed during the 2016 data taking period. The luminosity integrated over time is a measure of the amount of collected data. At $\sqrt{s} = 13 \text{ TeV}$, a total of 36.1 fb^{-1} of integrated luminosity was recorded in ATLAS.

3.2. The ATLAS detector

The ATLAS detector is a general purpose particle detection device. It is designed to offer a variety of subsystems capable of measuring particle interactions over a wide energy range. Several components are installed around the designated interaction point and serve different purposes. ATLAS consists of three major systems, the Inner Detector (ID), the calorimeter system, and the Muon Spectrometer (MS). The detector is divided into three sections along the beam axis. The *barrel* is the central part, at both ends of the detector two *endcaps* are placed too ensure high fiducial coverage. The overall layout of the detector can be seen in Figure 3.2. A coordinate system is defined where the y axis points to the surface and the x axis toward the center of the LHC ring. The z axis is set along the beam axis. Two angles are defined, ϕ is the azimuth angle in the xy plane. The angle θ describes the polar angle.

$$\phi \in [-\pi, \pi] \quad \theta \in [0, \pi] \quad (3.2)$$

$\phi = 0$ corresponds to the x direction. The *pseudo-rapidity* η is defined as

$$\eta = -\ln \left[\tan \left(\frac{\theta}{2} \right) \right]. \quad (3.3)$$

Pseudorapidity differences of two space points are Lorentz invariant under boosts along the beam axis, making it a convenient choice for analyses of hadron collider data. Momenta and energies of particles are commonly neglected in favor of their transverse components,

$$p_T = \sqrt{p_x^2 + p_y^2} \quad E_T = \sqrt{p_T^2 + m^2} \quad (3.4)$$

since these are invariant under Lorentz boosts along the beam axis, as well.

3.2.1. Measurement of particle tracks

The inner detector contains the tracking system of the detector. The tracking system serves two major purposes: it can determine particle momenta and allows for vertex reconstruction. The curvature

3. Experimental setup

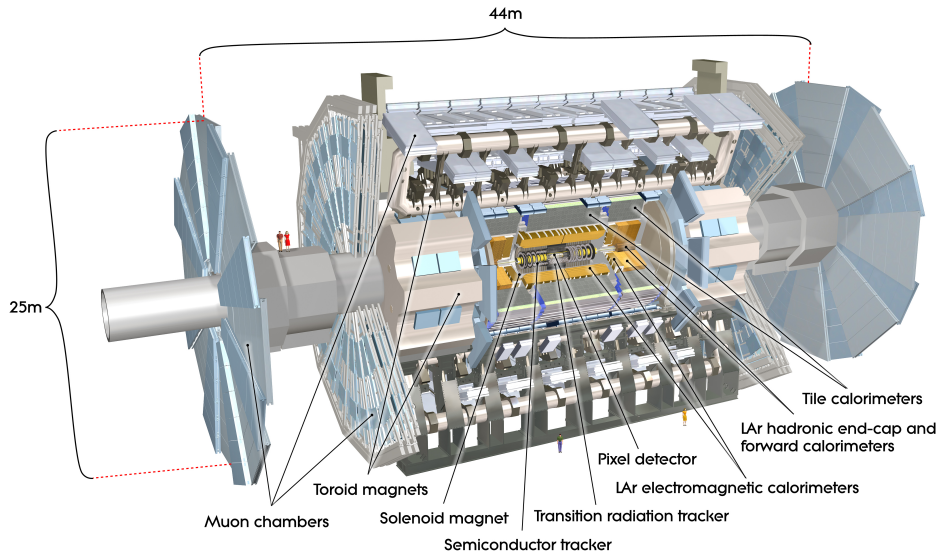


Figure 3.2.: Illustration of the ATLAS detector. It shows the various components used to detect particles created by the proton-proton collisions occurring in the beam pipe at the center. Figure taken from reference [25].

of charged particles, introduced by a magnetic field, allows for a momentum measurement. By extrapolating measured tracks to the interaction region at the center of the detector, interaction vertices can be reconstructed. One application of this is the identification of signatures belonging to pile-up vertices.

The following description is based on reference [26]. The tracking system consists of three subsystems. The first one is a silicon pixel tracker, followed by a silicon strip tracker called Semiconductor Tracker (SCT) and the Transition Radiation Tracker (TRT). The inner two components are semiconductor detectors. Particles passing the active segments create electron-hole pairs in the depletion zone created by a pn-junction between two semiconducting materials. This charge separation results in a detectable current in electrodes placed at opposing sides of the active material. By segmentation, spatial information can be gained. The pixel tracker is segmented in two dimensions, while the SCT is segmented in the form of strips. High granularity is required to improve the spatial resolution. Around the semiconductor trackers the TRT is installed. Transition radiation is produced when charged particles cross the boundaries of materials with varying dielectric constants. The TRT consists of straw tubes, which contain a gas mixture. In their center, a tungsten wire is placed. Particles traveling through the straw tubes ionize the gas mixture and a current can be measured in the anode wire. Material which facilitates the production of transition radiation surrounds the straw tubes. If a charged particle traverses this material, it creates transition radiation. The transition radiation photons can be detected by the straw tubes. The energy loss by transition radiation of a particle depends on the Lorentz factor γ . Consequently, particles of different masses produce different amounts of transition radiation. This assists the distinction of particle types. For instance, electrons and positrons produce more transition radiation than heavier charged pions. The large spatial extent of the TRT enhances track reconstruction by providing information at large distances from the interaction point.

Since the tracking detectors are close to the beamline, they are exposed to harsh radiation. The systems are designed to withstand radiation damage to a certain degree, a degradation in performance is expected regardless. An additional silicon pixel layer (Insertable B-layer (IBL) [27]) was installed as the innermost tracking layer during the long shutdown before the 2015 data taking run. This new layer counteracts the decreased resolution caused by degradation in the outer trackers and ensures good tracking performance. The additional tracking information also improves the identification of jets originating from b quarks, as well as mediating the effects of higher pile-up.

High spatial resolution is important for primary vertex reconstruction and analyses in general. The primary vertex is the vertex associated with the hard proton scattering and is defined as the vertex whose tracks have the highest $\sum p_T^2$. The design spatial resolution, is about $20\text{ }\mu\text{m}$ in the transverse plane and about $40\text{ }\mu\text{m}$ along the beam axis [26].

The tracking system stretches across the barrel region and the two endcaps. The pixel and SCT detectors cover a region of $|\eta| < 2.5$, while the TRT reaches up to $|\eta| < 2.0$. They are immersed in a magnetic field of 2 T created by a solenoid magnet. This field bends the trajectories of charged particles in the transverse plane, thus allowing for a measurement of their transverse momenta components. The relative momentum resolution scales with

$$\frac{\sigma_{p_T}}{p_T} \propto p_T. \quad (3.5)$$

This means that for higher momenta the resolution degrades because the curvature of the track decreases. For high momentum objects, such as the ones used in this analysis, the momentum measurement is thus worse than a calorimetric measurement of the energy (see Section 3.2.2). The sign of the charge of a particle can only be determined by considering the direction in which the track is bent. Since the curvature decreases, the separation between particles and antiparticles becomes inefficient for high momentum objects. Since the analysis is not sensitive to the sign of the lepton charge and to simplify the rest of the subsequent discussion, there will be no differentiation between particles and antiparticles.

Calorimetric energy measurement of muons is not feasible in the detector, due to them being minimally ionizing. Instead, only a measurement of their momenta is performed. To this end, a dedicated muon system surrounds the detector. This muon system consists of eight superconducting toroid magnets in the barrel and eight more in each of the endcaps. The magnets create an approximately circular magnetic field around the beam axis. Around these magnets, tracking detectors are placed. These additional tracking detectors can be used in combination with the tracking system in the ID to improve the muon momentum resolution. To be able to trigger a detector readout based on a muon signature, dedicated trigger chambers are also found here. The transverse momentum resolution is described by

$$\frac{\sigma_{p_T}}{p_T} = \frac{r_0}{p_T} \oplus r_1 \oplus r_2 p_T. \quad (3.6)$$

The first term marked by r_0 covers energy loss fluctuations along the trajectory. For muons, this occurs while traversing the calorimeter, whereas it is negligible for the track measurement of electrons in the ID. The second term marked by r_1 covers multiple scattering effects, as well as inhomogeneities of the magnetic field. The third term marked by r_2 covers resolution effects. For the muon

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system, this includes effects in the spectrometer, such as misalignment (see [28]). For electrons, r_0 vanishes, while $r_1 \approx 1\%$ and $r_2 \approx 0.05\%$. The momentum resolution achieved for muons is about $\sigma_{p_T}/p_T \approx 10\%$ at $p_T \approx 1$ TeV (values from [29]).

Figure 3.2 shows the various tracking system components, as well as the solenoid and toroid magnets.

3.2.2. Energy measurement

The detector has two major components which serve the purpose of measuring the energy of particles, the electromagnetic and the hadronic calorimeter. The following description is based on [26]. The calorimeter type used is a sampling calorimeter. These consist of alternating layers of active and inactive material. The inactive material is called absorber material. It serves the purpose of increasing the energy loss of passing particles and facilitating pair production of shower photons. This is desirable, since the calorimetric measurement requires the particle to deposit its entire energy in the device.

Particles passing the calorimeter interact with the material in a number of ways. Electrons, photons and hadrons create showers. The former two particles produce electromagnetic showers which are detected in the electromagnetic calorimeter. Electrons undergo bremsstrahlung, thereby producing photons. These high energy photons, as well as initially incoming photons, can create additional electron-positron pairs. This creation of additional particles happens repeatedly. The energy of the incident particle is distributed across all particles contained in the shower. It decreases as

$$E(x) = E_0 e^{-x/X_0}, \quad (3.7)$$

where E_0 is the initial energy, and X_0 is the radiation length. The radiation length is a material specific quantity describing the showering properties. When the shower particles reach the energy range where ionization dominates energy loss, no additional particles are created. The shower particles are detected in the active layers of the calorimeter. The energy of the incident particle is then proportional to the total number of detected particles.

Hadrons create a different type of shower. The particles undergo a sequence of inelastic collisions with the nuclei of the detector material. This produces additional hadronic particles which continue to shower. The absorption length λ is the equivalent hadronic quantity to the radiation length mentioned before. Since hadronic showers typically penetrate deeper into the material, the calorimeter size needs to be considerably larger to ensure full energy deposition within it. The passive absorber material also needs to be chosen carefully. During the scattering, hadrons also partially deposit energy without creating additional particles. Energy deposited in excitations of the nuclei and energy needed to overcome the nuclear binding energies is effectively extracted from the shower. This energy loss can be compensated in the calibration step. Another possibility is the use of an actively compensating calorimeter. By choosing an absorber material which captures the extracted energy and subsequently releases it back to the shower, a nearly full energy deposition can be achieved.

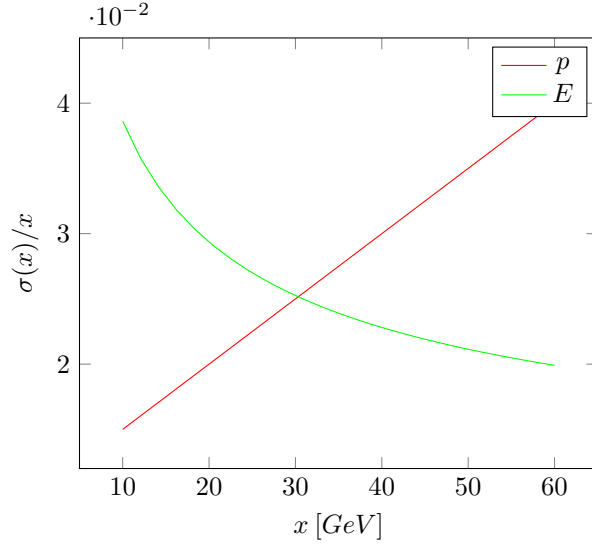


Figure 3.3.: Schematic comparison of the overall shape of relative momentum and energy resolutions.

Figure 3.3 shows the typical shape of the relative energy resolution of the calorimetric measurements of electrons, consisting of three terms:

$$\frac{\Delta E}{E} = \frac{a}{\sqrt{E}} \oplus \frac{b}{E} \oplus c \quad (3.8)$$

The calorimetric measurement outlined above corresponds to a counting experiment, the uncertainty on the counted number of particles is the Poisson error $\sqrt{N}$ which is proportional to $\sqrt{E}$. By design, a sampling calorimeter incorporates inactive material slices, which reduces the total resolution, while allowing for particles to be stopped within the calorimeter. These effects are described by the *sampling term* marked by a . Another aspect is the electronic noise which affects the actual readout of measurements from the instruments. This is taken into account in the *electronic noise term* marked by b . The term marked by c is the *constant term*, which increases the relative resolution independently of the energy of the particle. It originates from the absolute calibration of the calorimeter. In ATLAS, the sampling term factor is around $a \approx 10\%$ while the constant term amounts to $c \approx 0.7\%$ [29].

When comparing the energy resolution to the momentum resolution achieved by the tracking system, it is clear that for low energies or momenta, the energy measurement yields a worse resolution. Above a certain momentum, however, the energy resolution surpasses the momentum resolution. This can be seen in Figure 3.3, where the resolution lines intersect at $x \approx 30$ GeV. For higher momenta, the calorimetric energy measurement thus provides better resolution. This threshold depends on the involved components and on the region in η .

The electromagnetic calorimeter in ATLAS uses lead as an absorber and liquid argon as the active material. It covers the region $|\eta| < 3.2$ including components in both the barrel and the endcaps. For $|\eta| < 2.5$ the precision is markedly higher. The hadronic calorimeter is located behind the electro-

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magnetic calorimeter. In the barrel, the hadronic calorimeter uses iron absorbers and scintillator tiles as active material. In the endcaps liquid argon is used. It covers a region of $|\eta| < 3.2$ including barrel and endcap components. An additional forward calorimeter extending from $3.1 < |\eta| < 4.9$ delivers electromagnetic and hadronic energy measurements close to the beamline. By segmentation of the layers of the two calorimeter types, the position of energy deposits can be resolved. It also allows for a measurement of shower shapes. Angular information obtained from the calorimeters plays a crucial role in the reconstruction of particle candidates (see Section 3.2.4).

Figure 3.2 shows the electromagnetic and hadronic calorimeters in the barrel and the endcaps.

3.2.3. Event triggering and data acquisition

Due to the high instantaneous luminosity the rate at which collisions happen in the detector is high. At a proton bunch spacing of 25 ns, the bunch crossing frequency is 40 MHz. Multiple collisions can occur during each bunch crossing (see Section 4.4.3). Each collision is referred to as an event. The rate at which the detector data for events can be stored is much lower than the actual event rate. To mitigate this, a trigger system has to reduce the readout and storage rate. This can be done by buffering the detector readout, and evaluating it during operation. Only if an event is considered of interest, the data is actually read out and stored. The following is based on reference [30].

The trigger system in ATLAS consists of two levels, the higher level is only invoked if the lower level has already triggered. The hardware based lowest level one (L1) trigger looks at every single event seen by the detector. The L1 trigger combines information from the calorimeters and from the muon system. At this stage, the readout of the calorimeter occurs at a reduced granularity. By grouping the energy deposition in the calorimeter into clusters, the transverse energy of a signature can be calculated. In addition to this, a basic jet reconstruction can be performed, by summing the energy depositions from multiple clusters. This allows for additional requirements, such as on E_T^{miss} . The muon L1 trigger can identify patterns of muons in the muon tracking chambers. The information is fed into the Central Trigger Processor, where a combined trigger decision is performed. It can use the result of calculations carried out in a dedicated topological hardware trigger. The topological trigger can calculate complex geometric and kinematic properties and allows selection of events based on them. During the triggering procedure, care has to be taken to associate detector readouts with the correct bunch crossing. The L1 trigger system does not use any information from the tracking system in the ID.

The L1 reduces the total bunch crossing rate of 40 MHz to a trigger rate of about 100 kHz. For events accepted by the L1 trigger the detector readout is performed. The data is digitized and passed to the High Level Trigger (HLT) which runs on a server farm. If the L1 has identified specific regions of interest in the detector, the HLT is instructed to focus on these. At this stage, the information from the tracking system can be used in addition to the calorimeter and MS readout. For muons, the HLT makes use of dedicated Monitored Drift Tubes (MDT), which allow for a precise momentum

measurement. The HLT reconstructs objects seen in the detector and applies a selection on them. The reconstruction is performed in two passes, the first one with lower and the second with higher precision. This two-pass approach reduces processing time, by rejecting most of the events in the less expensive first pass.

The configuration of the trigger systems is described by a *menu*. The menu determines which requirements are set, and in which order. This results in chains of triggers. If a trigger chain has identified an event of interest, it is stored permanently. The information on which trigger chains have accepted an event is stored as well, allowing for offline selection of a specific trigger chain. The trigger menu is updated regularly to match the conditions in the detector. Since the bandwidth of the persistence system is limited, the trigger rate has to be kept within a certain boundary. As a consequence, the various thresholds need to be adjusted to meet the rate criteria, or a prescaling has to be introduced. A prescaled trigger only causes an actual readout for every N^{th} time it accepted an event, thus reducing its output rate. The trigger chains used in this analysis are listed in Section 5.1.

3.2.4. Particle reconstruction, identification and isolation

The data stored after the detector readout is a raw format. It contains all the information from the various components of the detector. This information needs to be converted into higher level physical objects, with properties such as energies and trajectories. This step is called reconstruction. The reconstruction algorithms are executed independently and perform pattern recognition on signatures contained in the full information from the detector. As a consequence, signatures in the calorimeter and tracks can be associated to more than one reconstructed object. These particle candidates are then tested for compatibility with the various particle hypotheses. Multiple identification levels are defined for each particle type. The levels correspond to varying strictness of the applied identification criteria.

To avoid considering particle signatures that overlap with other signatures, particles are tested for isolation. An example for an isolated electron can be found in the decay $W^- \rightarrow e^- \bar{\nu}_e$. The electron and neutrino are produced back-to-back. Hence, the electron will not be in vicinity of other signatures stemming from the same decay. The opposite scenario can for instance be found when considering a $t\bar{t}$ process. The top pair is produced back-to-back and both top quarks decay into a b quark and a W . The system of decay products is boosted due to the momentum of the tops. The W can decay electronically. This results in an electron signature close to the signature of the jet induced by the b quark. An isolation requirement can achieve separation between these two scenarios.

The isolation requirement is implemented by summing up the transverse energy from the calorimeter clusters, and the transverse momenta of tracks found in a certain cone around the particle in question. The size of the cone is fixed for calorimeter clusters and a function of the particle E_T for tracks. Clusters and tracks reconstructed as the particle signature itself are subtracted, to only take into account the adjacent activity. If the primary particle dominates the energy found in the cone,

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it is considered isolated. For this, the sums are divided by the transverse energy or momentum, respectively. Cuts are placed on the ratios whose specific values are exposed as working points.

This analysis and thus the following description focuses on electrons and muons. The isolation procedure is almost identical for electrons and muons, except for the specific choice of some variables.

Electrons

The following is based on [31]. For the reconstruction of electron candidates, energy depositions are identified as seeds in the electromagnetic calorimeter segments. From this seed a cluster of deposited energy in the calorimeter is build. In the tracking system, a track seed is created by combining hits from each of the three subsystems of the tracking system. An assumption on the expected energy loss as well as the impact of multiple scattering is made during the seed finding. Using this seed, a track is fitted to the hits in the trackers. Electron candidate tracks are required to be compatible with the primary vertex identified for the event. The tracks are then extrapolated and matched to the energy clusters in the calorimeter. If the track measurement is precise, but only loosely matches the energy cluster, the track is rebuilt using an electron specific approach which directly takes into account the possibility of bremsstrahlung. If a track and matching energy cluster is found, the particle is categorized as an electron. Calorimeter clusters without a track pointing to them are treated as one type of photon candidate, as these are not detected in the tracking system, but also produce a similar electromagnetic shower.

The electron candidate then undergoes an identification procedure. The algorithm considers parameters such as the calorimeter shower shape, information from the TRT, parameters related to the track-cluster matching step and other track quantities. Examples for shower shape variables would be the amount of the shower which penetrates the electromagnetic calorimeter and reaches the hadronic calorimeter, the energy deposit in the outermost layer of the electromagnetic calorimeter, and the shower width. Parameters related to the track include the number of hits in the various tracking detectors and the impact parameter of the track relative to the beamline. Track-cluster matching variables are the separation in η and ϕ , as well as the ratio between the measured energy and the measured momentum of the track. The aforementioned variables are evaluated using a likelihood approach. The probability density functions relative to electron *signal* and non-electron *background* hypotheses for the various parameters are used to calculate two likelihoods. A combined discriminant

$$d_{\mathcal{L}} = \frac{\mathcal{L}_S}{\mathcal{L}_S + \mathcal{L}_B} \quad (3.9)$$

is used to differentiate, where $\mathcal{L}_S$ and $\mathcal{L}_B$ are the likelihoods for the signal and background hypotheses, respectively. Additional cut requirements are imposed on the number of hits in various components. A set of three cuts is made on the discriminant $d_{\mathcal{L}}$ corresponding to three identification levels. These are typically referred to as "LHLoose", "LHMedium" and "LHTight". Tighter levels are subsets of the looser levels, and are optimized to meet certain signal efficiency requirements.

Muons

The following is based on reference [28]. Muon reconstruction is first performed separately for the inner detector and the MS. The ID track reconstruction is very similar to the electron equivalent described above. The reconstruction in the MS begins with hit patterns in the various chambers being recognized as *segments*. The tracks are fitted to the hits separately for the different MS subsystems, and then combined. The segments can be attached to multiple tracks, and are later either assigned to their best match or shared between tracks. Finally, the hits in assigned segments are fitted globally. Additional hits can be added at this stage when found compatible with the fit result. The track is then refitted using the newly added hits.

In a second step, tracks from the ID and MS are combined in a global refit using hits from both systems. Also at this stage, MS hits can be removed or added. The result is a *combined muon* candidate. There are complementary procedures, for instance the extrapolation of a track only found in the MS.

Identification of muons is done using several variables to assess the likely source of a combined muon. Among these are the difference of the charge-over-momentum ratio between ID and MS tracks, divided by the uncertainty, the difference between p_T of the ID and MS tracks, and the χ^2 of the combined fit. The identification is then divided into four levels. Three of them are "Loose", "Medium" and "Tight" with varying requirements on the aforementioned variables and the tracks themselves. "Loose" is optimized for reconstruction efficiency and allows all track combination options. "Medium" is intended to minimize the systematic uncertainties of the reconstruction, and uses combined tracks and tracks extrapolated from the MS. The "Tight" level is designed for purity, allowing only combined muons with high quality tracks. An additional "High- p_T " working point exists, which is optimized for muons with $p_T > 100$ GeV. For this, "Medium" level combined muons are used and additional requirements are imposed on them. The MS track needs to include at least three hits, and problematic regions in the MS are vetoed. As a consequence, the "High- p_T " working point provides the highest momentum resolution at very high momenta.

3.2.5. Calculation of Missing Transverse Momentum

As discussed in Section 2.3, the properties of the neutrinos in the signal final state need to be approximated via the missing transverse momentum. The vectorial sum of transverse momenta of all measured physical objects is calculated. The negative vector of this sum is then *missing* in order to restore conservation of momentum. In the case of an event producing a neutrino, this momentum can be attributed to the neutrino. Since signatures can be reconstructed as different particles, double counting has to be prevented by overlap removal. This can be achieved by requiring that particles are not within a cone of ΔR of each other. As described in [32], the terms taking into account different types of objects are

$$E_{x(y)}^{\text{miss}} = E_{x(y)}^{\text{miss},e} + E_{x(y)}^{\text{miss},\gamma} + E_{x(y)}^{\text{miss},\tau} + E_{x(y)}^{\text{miss},\text{jets}} + E_{x(y)}^{\text{miss},\mu} + E_{x(y)}^{\text{miss},\text{soft}}. \quad (3.10)$$

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The final value for E_T^{miss} is then simply

$$E_T^{\text{miss}} = \sqrt{(E_x^{\text{miss}})^2 + (E_y^{\text{miss}})^2} \quad (3.11)$$

and the angle ϕ in the transverse plane can be calculated as

$$\phi^{\text{miss}} = \arctan(E_y^{\text{miss}}/E_x^{\text{miss}}) \quad (3.12)$$

The sum is split in terms stemming from electrons, photons, taus and jets. The selection of these objects follows the optimized prescription used for the published analysis. Any tracks that are not associated with these objects is accounted for in the track soft term, contributing as $E_{x(y)}^{\text{miss, soft}}$.

It is obvious that detector effects can severely impact the computation of this value. For instance, any particle whose momentum is either miss measured within, or escapes the acceptance of the detector, shows up as *missing*. It is also important to use calibrated objects in the reconstruction, so as to get a correct magnitude for E_T^{miss} .

The calculation of E_T^{miss} is implemented in a centrally maintained software package (see [33]). The software receives particles explicitly selected and calibrated at the analysis level and handles the calculation automatically.

3.2.6. Luminosity determination

As described in Chapter 2, the luminosity is an important ingredient to connect cross-sections to the number of observed events. Therefore, it is crucial to obtain a precise luminosity measurement during the operation of the detector. Since the luminosity affects all processes equally, it can be inferred from the total activity in the detector. The following description of the luminosity measurement methodology is based on reference [34].

The total activity in the detector includes all particles produced in the inelastic proton-proton scattering processes that take place during a bunch crossing. The bunch luminosity can be calculated as

$$\mathcal{L}_b = \frac{\mu_{\text{vis}} f}{\sigma_{\text{vis}}} \quad (3.13)$$

where μ_{vis} is the visible interaction rate per bunch crossing, σ_{vis} is the visible cross-section and f is the revolution frequency. μ_{vis} can be measured during operation. There are three independent measurement approaches, two of which use dedicated detectors. The BCM⁴ consists of diamond sensors placed approximately 2 m on either side of the nominal interaction point. The detector reads every sensor activation as a hit and is able to be read out more than once per nanosecond. The ATLAS luminosity monitor LUCID⁵ uses Cherenkov radiation created by passing particles. LUCID consists of two detectors, positioned along the beam axis on either side, 17 m away from the interaction point. The Cherenkov detectors are monitored using photomultiplier tubes. Both BCM and LUCID count

⁴Beam Conditions Monitor

⁵Luminosity measurement using a Cherenkov Integrating Detector

hits for each bunch crossing. The bunch crossings are labeled sequentially by an identifier (BCID). This way, the number of hits for each specific BCID can be monitored across an entire proton fill of the LHC. In both cases, the number of hits can be converted into the visible interaction rate μ_{vis} .

Another possibility is to use the tracking system of the detector itself. The number of reconstructed tracks is a measure of the interaction rate. Since the readout of the tracking system cannot be performed for every bunch crossing, the number of tracks is integrated over a *luminosity block*, typically around one minute. To not be biased by a specific trigger setup, a random trigger with a frequency of about 100 Hz is used.

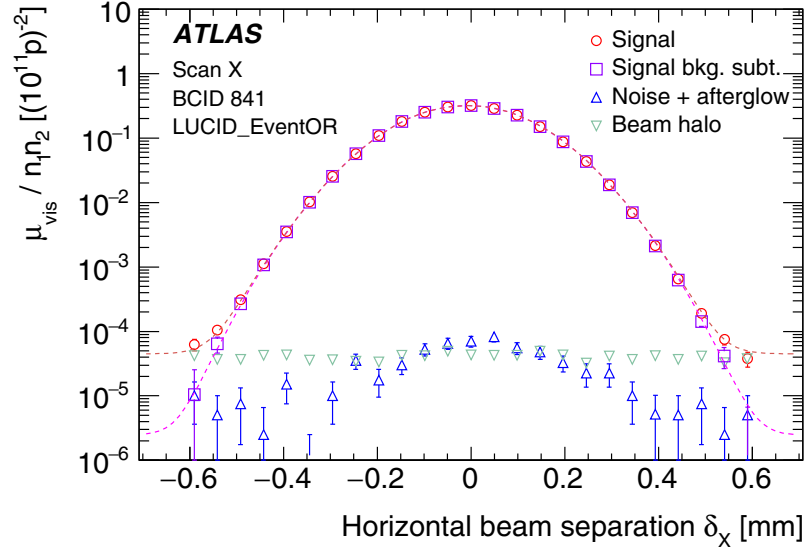


Figure 3.4.: Visible interaction rate depending on the separation of the beams during a vdM scan. Figure taken from [34].

The measurement from either of these methods requires a calibration to obtain the actual luminosity value. This absolute luminosity measurement is done in dedicated runs in advance of data taking. In a *vdM* scan, the two incoming beams are shifted in opposite directions before entering the interaction region. During runs dedicated to the vdM scans, the bunch intensity is lowered, and only a few bunches are injected into the collider. These conditions reduce pile-up and help reduce the systematic uncertainties arising from the calibration procedure. The shift results in an offset between the incoming beams. This offset causes a decrease in the observed luminosity. The shift is performed subsequently in the x or y direction. The bunch luminosity can then be calculated as

$$\mathcal{L}_b = \frac{f n_1 n_2}{2\pi \Sigma_x \Sigma_y} \quad (3.14)$$

where n_1 and n_2 are the number of protons in each bunch per beam. Σ_x , Σ_y are a measure of the size of the overlap of the beams in x and y . The beams do not have to be Gaussian. Σ_x and Σ_y can be calculated as

$$\Sigma_i = \frac{1}{2\pi} \frac{\int R_i(\delta) d\delta}{R_i(0)} \quad (3.15)$$

under the assumption that they factorize. A similar calculation can also be performed without that assumption. $R_i(\delta)$ is the luminosity at a beam separation of δ . A plot showing the visible interaction

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rate for one BCID as a function of the horizontal beam separation during a vdM scan is shown in Figure 3.4. The absolute luminosity can be extracted by evaluating the width of the distribution shown here. In the case of a Gaussian beam, the width parameters Σ_i correspond to the standard deviation of the distribution. Using the value for $\mathcal{L}_b$ obtained from a vdM scan and a simultaneous measurement of the BCM, LUCID or the tracker, the corresponding visible cross-section σ_{vis} can be calculated, and used as a calibration for the online luminosity measurement.

The LUCID system is used for the actual measurement of the luminosity during 2015 and 2016 data taking. The results obtained using the three approaches are compared and tested for consistency. The luminosity is integrated over time for a series of luminosity blocks. The detector conditions are unchanged during each luminosity block. The integrated luminosity is a measurement of total amount of collected data. Figure 3.5 shows the integrated luminosity over time for the 2015 and 2016 data taking periods. Both the effective luminosity delivered by the LHC and the actual luminosity which was recorded in the ATLAS detector are displayed. The 2015 integrated luminosity rises steeply after an initial slower rise. In 2016 the integrated luminosity rises with an approximately constant slope during the entire online period. For this analysis only luminosity blocks during which the detector was operating nominally without failures are used. The integrated luminosity and a list of viable luminosity blocks is provided centrally. The uncertainty on the total integrated luminosity of the data is calculated centrally to be 3.2 %. It is considered as a systematic uncertainty in this analysis.

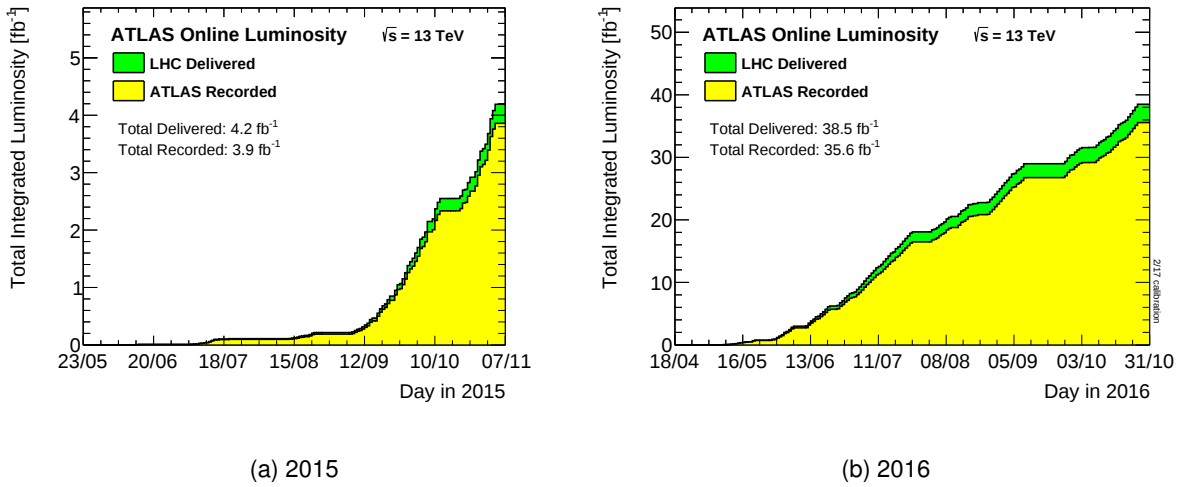


Figure 3.5.: Integrated luminosity for different days in 2015 and 2016. The integrated luminosity rises monotonously as more data is recorded. Figures from [35].

4. Simulation

When working to analyze any kind of collected data, implicit or explicit assumptions and expectations are relevant to the interpretation. In physics, many processes are understood at a theoretical level, which allows for their simulation. Simulated results of physics experiments can be used to characterize devices and instruments, as well as help test and develop both hardware and software. The simulation of physics processes and the detector are centrally conducted by the ATLAS collaboration, and the corrections and calibrations which are applied at the analysis level are provided centrally.

4.1. Monte Carlo Simulations

One very common approach to simulation is called Monte Carlo (MC) simulation. Monte Carlo methods make use of random numbers to calculate the predictions. An example of a Monte Carlo method, albeit not a physics simulation, is integration. For instance, calculating the value of π is straight forward using uniformly distributed random numbers as well as Pythagorean geometry.

A crucial input to all kinds of Monte Carlo methods is high quality random numbers, meaning true uniform distributions, from which arbitrary other distributions can be easily sampled. One approach is to use actual random numbers, taking for instance a temperature sensor readout which measures and samples truly random fluctuations in ambient temperature. Another approach, and one that is predominantly adopted for Monte Carlo methods, is pseudo-random numbers. Pseudo-random numbers are sequences, generated via algorithmic calculation and have well defined properties. With a carefully built algorithm using an appropriate choice of input parameters, some of which need to be prime numbers, these pseudo-random numbers can achieve very high uniformities. An added benefit is that these sequences are reproducible. Using the same seed value as a starting point, a pseudo-random number generator always yields the exact same sequence. This is useful in the development of algorithms using these random numbers. If their input is reproducible, re-running it does not result in a different outcome.

The simulations that are required for generating background and signal expectations which can be used in an analysis can be divided into two parts: the simulation of the underlying physics processes themselves and the simulation of the response which the detector generates if such a process occurs within it.

4. Simulation

4.1.1. Simulation of physics processes

Using Monte Carlo techniques to simulate physics processes at the fundamental level works by approximating the non-deterministic nature of quantum physics, and quantum field theories in particular. Locations, momenta and all processes are attached to probability densities, resulting in the production of a specific outcome. The challenge for simulation is therefore to model these probability densities, and sample from them to achieve a result which resembles reality.

The initial process simulation uses the analytical cross-section which can be determined from the EW theory or QCD. Since these theories rely on perturbation theory to approximate the result of infinite sums of terms, a cut-off order of the coupling constants α , α_s has to be chosen. For the simulated samples used in the W' analysis, background simulations are performed at next-to-leading-order (NLO). There are uncertainties associated with the coupling constants used in the generators, which affect the analysis systematically. Another input that is required to determine the cross-sections is the parton distribution function, of which multiple applicable ones are available. Differences arising from the choice of a PDF, as well as the uncertainty on the PDF itself can be accounted for as systematic uncertainties.

Using weighting factors that can be determined in independent and non-trivial procedures, the NLO simulation can be taken to higher orders, in particular to NNLO. This process is a phenomenological one, which takes the outputs generated at different orders, and uses their ratio to determine a weighting factor. These so called *k-factors* introduce another systematic uncertainty which needs to be propagated to wherever the reweighted sample is used.

Besides the simulation of the process of interest on the parton level, several subsequent steps have to be performed. Having been stripped of one of their partons, the proton remnants can undergo additional reactions which might result in particles which show up in the active volume of the detector. In a lot of cases, the majority of this effect is suppressed by the fact that the proton remnants typically have low transverse momenta. This means that they exit the detector along the beam axis. The involved partons, as well as the decay products, can undergo radiation before and after the primary interaction. As a consequence additional particles can be produced which may enter the active volume. Another process which needs to be treated is hadronization, where color charged particles, that are left as free particles after interaction, form color-neutral particles by producing particle-antiparticle pairs. Note that all of these processes are part of the initial event being simulated, and do not result from interaction with the detector material, which gets handled in a separate step.

Information about the particles at this generator stage is stored in the resulting sample, and is called the "truth" record. This record contains the true particle type, the true energy and momentum, and the true track of the particle. At truth level, the connections between particles are also accessible.

4.1.2. Simulation of the detector response

After having successfully produced a set of particles and momenta which correspond to the products of a process being simulated, the next step is to approximate how the detector reacts to them. To this end, the propagation of the produced particles through the detector material is simulated. Here, the different modes of interaction and their probabilities have to be considered. Additionally, a geometry model of the detector, putting all components in the right places, as well as accounting for the respective materials is needed. In high energy physics in general, and in ATLAS in particular, the software package GEANT4 [36] is used for this.

GEANT4 simulates the trajectories of particles in steps, taking into account momentum losses and direction changes due to electromagnetic and hadronic interactions as the particles move through the surrounding material. The software also handles secondary particles and their interactions. In the end, the quantity of interest is the deposited energy in each active component of the detector. This energy deposition can then be converted into a simulated signal which closely resembles that of an actual readout of the detector. The signal can then be fed into the reconstruction algorithm, just as an actual readout from the detector is. As a final result, the so called reconstruction level contains the same information, that would be available for actual data events. The truth record is stored alongside the information at reconstruction level for each event. This enables study of the reconstruction result, while still retaining access to the truth information when necessary.

4.2. Simulation of background samples

A problem that is encountered when trying to use Monte Carlo samples in an analysis is statistics. Ideally, the simulated sample has high statistics across the entire range that is relevant to the analysis. When simulating a physical process, off-resonance production of particles is suppressed, the suppression typically increases towards higher invariant masses. This problem can be mitigated by producing the sample in mass bins, where each bin contains a fixed number of events with invariant masses falling inside the bounds. To arrive at the final distribution, the distributions for the various mass bins are scaled to a specific integrated luminosity, using the effective cross-section of the process inside the bin and the total number of events. After this, the joint distribution can be treated like an inclusive one. The samples for the charged and neutral current Drell-Yan processes are mass binned, while the other samples are inclusive. The mass binned samples are provided centrally, while the scaling and combination is performed at the analysis level.

An image which visualizes this procedure can be seen in Figure 4.1. It shows the invariant mass distribution for the decay $W^+ \rightarrow e\nu$ at truth level. The various mass-bins are drawn in different colors. The edges of the bins match precisely, the entire distribution fits together well. Note that in order to see this image, the variable on which the mass-bins are defined has to be studied. The corresponding plot of different kinematic variables do not show clear separation, as shown in Figure 4.2.

4. Simulation

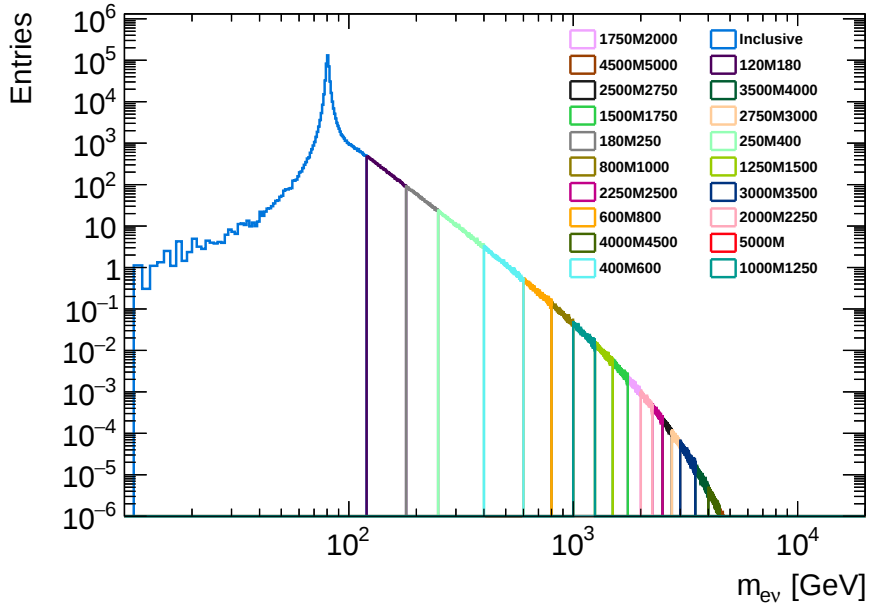


Figure 4.1.: Distribution of the invariant mass $m_{e\nu}$ for the decay $W^+ \rightarrow e^+ \nu_e$ at truth level from different mass bins.

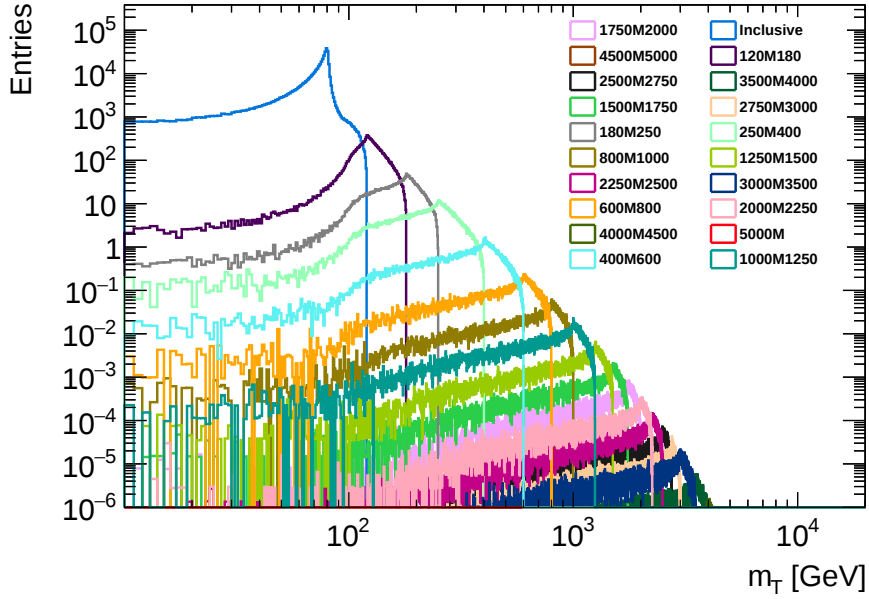


Figure 4.2.: Distribution of the transverse mass for the decay $W^+ \rightarrow e^+ \nu_e$ at truth level, from different mass bins.

Table 4.3 shows the four categories of backgrounds, W , Z^0/γ^* , Top and Diboson. The table also contains the specific processes that are included in each category. In order to generate the background samples for the charged and neutral current Drell-Yan processes as well as the Top and Diboson backgrounds, different generator setups are used. Details can be found in Table 4.2.

The Drell-Yan backgrounds are produced using POWHEG+PYTHIA8 [37, 38], while the Diboson backgrounds are simulated with SHERPA [39, 40]. Final state radiation of the decay products is added to the simulated events using PHOTOS [41] for the Drell-Yan processes. These backgrounds are corrected to NNLO in QCD and NLO in EW by applying mass dependent k -factors. The k -factors are calculated in a procedure outside of the scope of the analysis conducted for this thesis (see references [42, 43]). For the QCD corrections, factors of 1.03 (1.04), 1.02 (1.01) and 1.09 (0.87) for W^+ (W^-) at 0.5 TeV, 2 TeV and 4 TeV are applied. The EW corrections amount to 0.95 (0.95), 0.86 (0.86) and 0.81 (0.80) for W^+ (W^-) at 0.5 TeV, 2 TeV and 4 TeV. For the neutral current processes, QCD factors of 1.04, 1.02 and 0.94 at 0.5 TeV, 2 TeV and 4 TeV are used, whereas for the EW corrections the values are 0.99, 0.92 and 0.88 [29].

The $t\bar{t}$ and single- t backgrounds are generated using POWHEG [37]. PHOTOS is used for these backgrounds as well. All background samples are simulated at NLO using the CT10 NLO [44] parton distribution function.

Table 4.1 shows the various processes that are simulated, as well as their cross-sections and integrated luminosities. The cross-sections for the dominant W^+ and W^- processes are highest at 11.35 nb and 8.31 nb, respectively. The neutral current processes all have a cross section of 1.92 nb. The Diboson and Top background process groups have a combined cross-section of < 1 nb.

Table 4.1.: The different Monte Carlo samples and their corresponding number of events and cross-sections. For the Z and W processes, events in the inclusive sample which overlap with the mass bin samples are discarded.

sample	N_{evt}	σB [nb]	$\mathcal{L}_{\text{int}}$ [fb $^{-1}$]
$W^+ \rightarrow e^+ \nu_e$	31.60×10^6	11.35	2.79
$W^- \rightarrow e^- \bar{\nu}_e$	21.65×10^6	8.31	2.61
$W^+ \rightarrow \mu^+ \nu_\mu$	31.64×10^6	11.35	2.79
$W^- \rightarrow \mu^- \bar{\nu}_\mu$	21.70×10^6	8.31	2.61
$W^+ \rightarrow \tau^+ \nu_\tau$	31.58×10^6	11.35	2.78
$W^- \rightarrow \tau^- \bar{\nu}_\tau$	21.70×10^6	8.31	2.61
$Z^0 \rightarrow e^+ e^-$	21.80×10^6	1.92	11.34
$Z^0 \rightarrow \mu^+ \mu^-$	39.42×10^6	1.92	20.50
$Z^0 \rightarrow \tau^+ \tau^-$	31.32×10^6	1.92	16.29
Diboson	20.89×10^6	0.12	168.22
Top	69.33×10^6	0.83	83.17

4. Simulation

Table 4.2.: Table of the generators and parton distribution functions used to produce the various samples.

sample	generator	order at gen. level	PDF	corrected to order
$W^\pm \rightarrow l\nu$	POWHEG + PYTHIA8 [37][38]	NLO (QCD+EW)	CT10 NLO [44]	NNLO QCD, NLO EW
$Z^0/\gamma^* \rightarrow ll$	POWHEG + PYTHIA8 [37][38]	NLO (QCD+EW)	CT10 NLO [44]	NNLO QCD, NLO EW
WW, ZZ, WZ	SHERPA v2.1.1.4 [39]	NLO QCD, LO EW	CT10 NLO [44]	-
$t\bar{t}$, single- t	POWHEG [37]	NLO QCD, LO EW	CT10 NLO [44]	-
$W' \rightarrow l\nu$	PYTHIA8 [38]	LO (QCD+EW)	NNPDF23 LO [37]	NNLO QCD, LO EW

Table 4.3.: Decay processes which are considered as background sources.

W	Z^0/γ^*	Diboson	Top
$W^\pm \rightarrow e\nu$	$Z^0 \rightarrow e^+e^-$	$Z^0 Z^0 \rightarrow llll$	$t\bar{t} \rightarrow lX$
$W^\pm \rightarrow \mu\nu$	$Z^0 \rightarrow \mu^+\mu^-$	$W^\pm Z^0 \rightarrow lll\nu$	$t \rightarrow lX$
$W^\pm \rightarrow \tau\nu$	$Z^0 \rightarrow \tau^+\tau^-$	$W^\pm W^\pm \rightarrow ll\nu\nu$	$\bar{t} \rightarrow lX$
		$Z^0 Z^0 \rightarrow ll\nu\nu$	$W^\pm t$
		$W^\pm Z^0 \rightarrow l\nu\nu\nu$	$W^\pm \bar{t}$
		$W^+ W^- \rightarrow l\nu qq$	
		$W^+ W^- \rightarrow qq l\nu$	
		$W^\pm Z^0 \rightarrow l\nu qq$	
		$W^\pm Z^0 \rightarrow qq ll$	
		$Z^0 Z^0 \rightarrow qq ll$	

4.3. Simulation of W' signal samples

In addition to the samples used for the background, another set of samples is needed. The data obtained is tested against a series of pole mass W' hypotheses, each pole mass yields a different spectrum, where the resonance position reflects the value of the pole mass. The obvious approach would be to simulate the signal hypothesis at every desired pole mass. The simulation, however, needs to include both the signal process, as well as the detector response. The latter is highly expensive computationally. Thus, simulation of a high number of pole masses would be prohibitively time consuming. Another problem is that very high numbers of events would have to be generated to ensure high statistics across the required mass range, or a mass binned approach, as it is used for some of the background samples, would have to be adopted. Instead, another approach is chosen.

By removing the Breit-Wigner term of the cross-section, and modifying the cross-section by multiplying a function of the invariant mass $m_{e\nu}$, a sample can be generated which is practically flat in mass. This means that the sample has almost equally high statistics for all values of the parameter of interest. The flat sample is generated using PYTHIA8 [38] at leading-order (LO). The distribution of m_T for the electron channel, calculated at truth level, can be seen in Figure 4.3. Final state radiation is included using PHOTOS at the generation stage. The sample receives corrections to NNLO

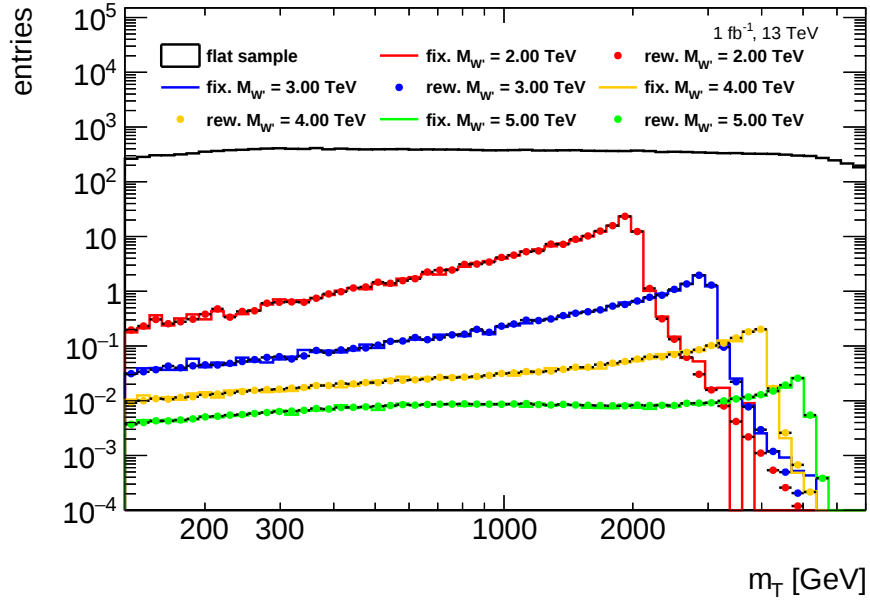


Figure 4.3.: Comparison of m_T distributions of the flat signal sample, results of the reweighting procedure and fixed-mass simulations. The distributions are scaled to an arbitrary integrated luminosity of 1 fb^{-1} .

in QCD, by application of a mass dependent k -factor. This weighting introduces drops at the low and high mass edges of the spectrum, which are not visible in the uncorrected sample. Of course, this flat distribution is unphysical. This means it cannot immediately be used as a signal expectation. Using a weight which depends on the pole mass and the invariant mass $m_{l\nu_l}$, the flat sample can be reweighted to the expected resonance shape. Thereby, the mass dependence is effectively reintroduced. The definition of the weight can be found in Chapter 6. To make sure that the result of the reweighting procedure yields reasonable shapes, full simulation of separate samples is performed at fixed pole masses:

$$M_{W'} \in \{2, 3, 4, 5\} \text{ TeV} \quad (4.1)$$

The samples are then compared to the result of the reweighted flat sample at corresponding pole masses. The comparison can be found in Figure 4.3. Plotted in color, both types of samples are shown, the dots belong to the distributions obtained from the flat sample, while the solid line of the same color shows the corresponding separately simulated sample. The distributions for the various masses rise until they reach their respective pole masses, and drop steeply for higher masses. Ratios of the distributions can be seen in Figure 4.4 on the following page.

From this comparison, it can be seen that the reweighting procedure yields satisfactory results. The reweighted distributions appear smooth across the mass range, since they contain the same statistics as the flat sample.

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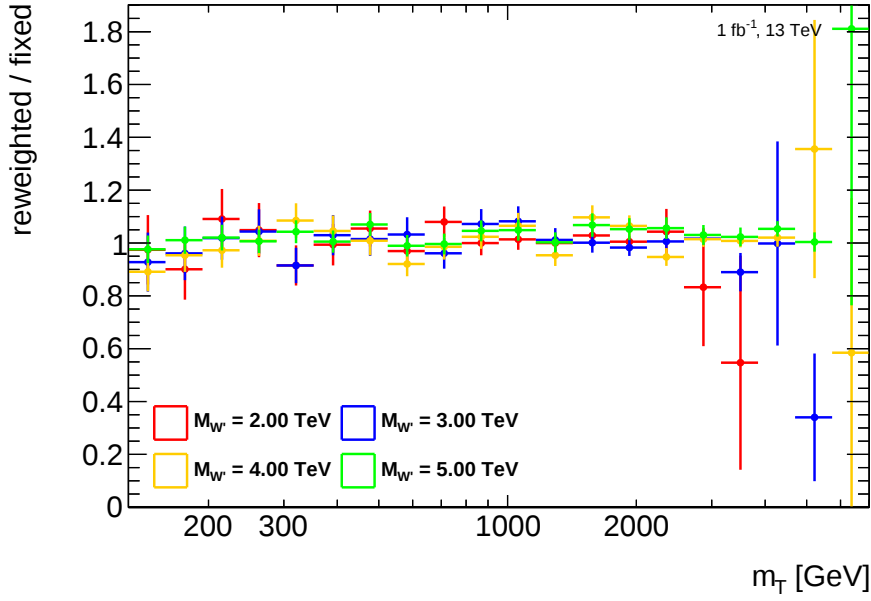


Figure 4.4.: Ratio of m_T distributions of reweighting results and fixed-mass simulations, at various pole masses.

4.4. Empirical corrections of simulated samples

Even though the Monte Carlo simulations reproduce the underlying physical processes as well as the geometry and material of the detector, differences between the distributions can still arise. The modeling of the aforementioned properties might not be perfect, or in some cases, exceed precision achievable in actual operation. To account for this, methods exist to match up the prediction with the actual measurement. This analysis focuses on the electron channel, where all recommended corrections of the simulated samples are applied. While this is also recommended for the muon samples, the correction factors are not applied, as a way to simplify the analysis.

4.4.1. Energy calibration

The energy measurement which is performed by the detector must be calibrated. The first step is hardware calibration, which is necessary to produce energy values in sensible units. However, discrepancies between actual data and expectation from simulation might still be present.

A correction of the energy calibration of electrons is done in a sequence of steps. Initially, a simulation based energy calibration is applied to both simulated samples and data, bringing them to a first agreement. Corrections are then applied to data events, which mitigate any non-uniformities that might be contained in the detector response. $Z^0 \rightarrow e^+e^-$ decays are then used as a reference to determine a set of two corrections. On one hand, the energy scale offset is determined, and

the data events are shifted so that they match the expectation. This is done in data, so that other energy measurements reproduce expected values. On the other hand, the energy resolution found in simulation is smeared to reproduce the actual resolution found in data. Details of the procedure can be found in [45, 46].

Since the calculation of the E_T^{miss} , which is described in Section 3.2.5, includes jets, the calibration of these objects is also applied in this analysis. Jets produce hadronic showers. Therefore, calibration needs to account for energy losses not measured by the calorimeter (see Section 3.2.2).

For muons, a similar procedure (see [28]) is used to correct the momentum scale in data as well as the momentum resolution in simulation.

4.4.2. Efficiency corrections

Another aspect which shows slight deviations in data and simulation are efficiencies of the selection criteria. The trigger, isolation, reconstruction and identification (see Section 3.2.4) procedures result in different efficiencies. These efficiencies can be evaluated for electrons and muons using a tag-and-probe approach [31]. This works by identifying pairs of leptons from the decays $Z^0 \rightarrow l^+l^-$ and $J/\Psi \rightarrow l^+l^-$. One of the leptons is required to produce a signal in the detector, and meet strict criteria. Since the leptons are produced in pairs and originate from the same decay, the missing lepton can be found as well, without using the requirement under study. By assessing the probability of this other lepton to meet the requirement, the efficiency of the criterion can be determined. By applying the same algorithm on data and simulated events, the efficiencies for both cases can be measured. A scale factor can subsequently be derived, which is used to correct the simulation to match the efficiencies seen in data.

The electron identification efficiency was found to be above 0.8, while the isolation efficiency was found to be around 1.0 after turn-on. The trigger efficiency was found to be close to 0.95 after turn-on and reconstruction efficiency was found to be around 0.98 (values from [31]).

For muons, the reconstruction efficiency was found to be around 0.9 for the High- p_T working point, the isolation efficiency was determined to be above 0.9, depending on the working point. The identification efficiency was measured to be around 0.8 for the High- p_T working point (values from [28]).

4.4.3. Pile-up effects

As mentioned in Section 3.1, bunches of protons collide in the detector every 25 ns. The instantaneous luminosity (see Equation (3.1)) seen by the detector depends linearly on the number of protons per bunch. To increase the rate at which rare processes take place, an increase of the luminosity is sensible. A side effect of higher instantaneous luminosity is the so-called *in-time* pile-up. Additionally,

4. Simulation

the bunch crossing rate influences the impact of particle signatures from previous bunch crossings, called *out-of-time* pile-up.

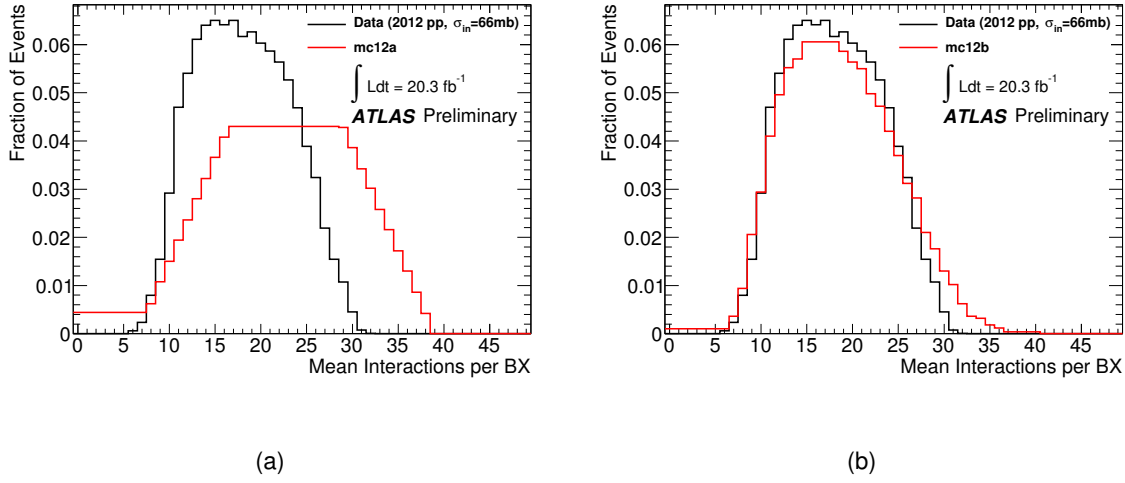


Figure 4.5.: Two plots showing the distribution of the mean number of interactions $\langle\mu\rangle$ in data and MC. The plot on the left shows the MC production run *mc12a*, which features a flat topped $\langle\mu\rangle$ distribution, while the plot on the right shows *mc12b*, whose $\langle\mu\rangle$ distribution roughly matches the one found in data, event before reweighting. Figures taken from [47]

The number of interactions per bunch crossing is called μ , and the value which is often discussed is the mean number of interactions $\langle\mu\rangle$ for a given period of data taking. Pile-up interactions do not necessarily have to be of the same type as the interaction which triggers an event in the detector, and do not have to be located in vicinity to the primary vertex. One possibility to suppress the influence on pile-up in the measurement is in fact to require particle tracks that make up the event to be within short distance of the primary vertex. Nevertheless, the effect of pile-up has to be modeled in simulation. To achieve this, the simulation produces events with different pile-up conditions, meaning a specific distribution of $\langle\mu\rangle$. Particles coming from pile-up events can influence the identification of particles. For instance, the shape of the electromagnetic shower of an electron, which is used in the identification, can be affected. The identification procedure includes measures intended on mitigating the impact of pile-up. This necessitates a good modelling of pile-up and primary vertex multiplicities in simulation.

These distributions can be seen in Figure 4.5 in red. Here, two sets of distributions are shown, the one in Figure 4.5a shows the pile-up conditions simulated before data taking at $\sqrt{s} = 8$ TeV began in 2012. To be conservative, a flat distribution, which spans a wider range of $\langle\mu\rangle$, was chosen. After data has been collected, the corresponding distribution of $\langle\mu\rangle$ can be determined for the data set, which is shown in black. Using these distributions, a weight can be derived which harmonizes the pile-up condition found in MC to the actual one in data.

It is apparent, that the flat distribution in MC necessitates that some events in the flanks receive a large weight, which leads to high statistical uncertainties. To remedy this, a second simulation is run, which already roughly reproduces the shape of the $\langle\mu\rangle$ distribution of the data set it is supposed to

model, as is shown in Figure 4.5b. Using the same reweighting approach, the distributions can again be brought to agreement, resulting in better utilization of the available statistics in MC.

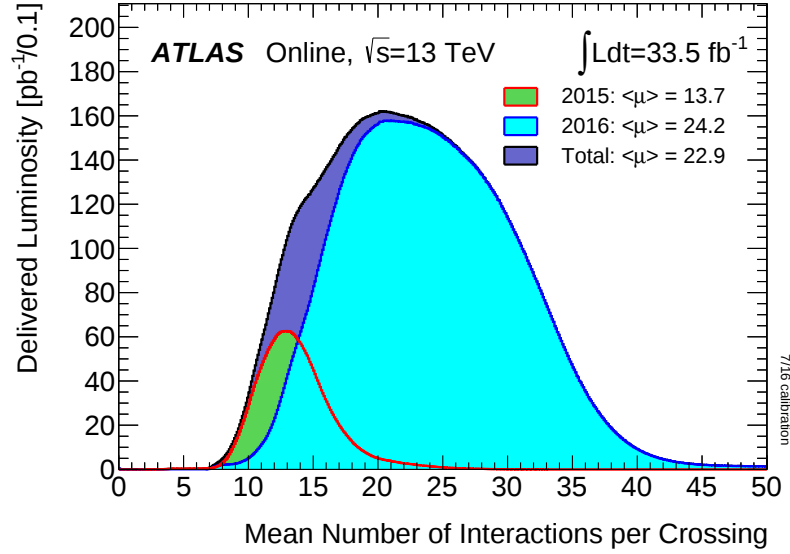


Figure 4.6.: Distribution of $\langle\mu\rangle$ for the data taking periods in 2015 and 2016. The figure is taken from [35]

Pile-up conditions have changed since the 2012 data taking period, since instantaneous luminosity has been increased substantially for 2016. For the first round of simulation that was done for $\sqrt{s} = 13$ TeV, a $\langle\mu\rangle$ distribution with a flat top was used again. The flat part starts at around $\langle\mu\rangle \approx 10$ and ends at $\langle\mu\rangle \approx 30$. Subsequent simulation rounds used $\langle\mu\rangle$ distributions determined from data.

Figure 4.6 shows the distributions of $\langle\mu\rangle$ seen in data for the 2015 and 2016 data taking periods. The distribution for 2015 has a maximum at around $\langle\mu\rangle \approx 13$, which is below 2012 values. On the other hand, the 2016 distribution peaks at $\langle\mu\rangle \approx 24$, clearly above previous values. The combined distribution for 2015 and 2016 is also shown. The increased instantaneous luminosity and resulting mean number of interactions per bunch crossing poses challenges for all stages of analyses. However, steps have been taken in all referenced efforts to cope with these new pile-up conditions.

The monitoring of pile-up conditions and the development of the reweighting procedure described above are handled centrally by the ATLAS collaboration. The reweighting is applied at the analysis level.

5. Analysis procedure

A search for W' bosons as described by the SSM is performed. Centrally provided containers are used as inputs for the data and the simulated samples. These containers have undergone the processing described in Chapter 3. They are read by an algorithm specifically created for this analysis. The algorithm evaluates every event, and makes use of centrally developed software packages to handle calibrations, corrections and certain calculations that have to be performed at the analysis level. It accesses the properties of the particles found in the containers at reconstruction level. At this level, the data and simulated samples can be treated equivalently. In the case of simulated samples, the algorithm can also access the truth record, that contains information on the particles simulated by the Monte Carlo generator.

The following chapter on the procedure begins with a description of the selection that is made on the events. A data driven approach to determine the Multijet background is introduced. The kinematic properties of the events passing the selection requirements are discussed, as well as the systematic uncertainties which affect the background estimates. A background extrapolation is used to remedy the low statistics at high transverse masses in the Top, Diboson and Multijet backgrounds.

5.1. Event selection

After the collection of data, the types of events which are to be studied can be selected using a variety of different criteria. These criteria need to be chosen in a way which selects the desired signature, while at the same time suppressing background processes, which feature similar signatures. Careful studies of the effects of the application of additional criteria have to be performed, in order to arrive at an optimal set. These optimizations have been conducted during previous ATLAS analyses such as [48, 49, 50, 51, 52], and this analysis makes use of a well established set of selection criteria, shown to produce desirable results. Even after tuning of selection criteria, background processes, which feature the same final state as the signal process under study, remain. These backgrounds are thus called irreducible. In the search for the decay $W' \rightarrow l\nu$ the events need to contain either an electron or a muon. To simplify nomenclature, the term lepton will cover both muons and electrons in the following, but exclude taus, since those are not considered in the analysis.

Trigger

The first requirement is that the event has been recorded due to a corresponding trigger having selected a lepton candidate. The specific triggers which are used are found in Table 5.1. The used triggers come with a certain energy threshold, and a mixture of triggers using different identification criteria is used for reasons explained in Section 5.2.2 on the data driven Multijet background estimation. The efficiency, with which the triggers select events, depend on their requirement value of p_T . Since the search is for resonances in the high-mass end of the spectrum, only high- p_T leptons are relevant. Requirements on the value of p_T from an off-line calculation are imposed. The precise values which are required are chosen so that the efficiency curve for the lowest trigger has already reached a very high level. The cut value on the off-line p_T for electrons is therefore $p_T > 65$ GeV, while the one for muons is at $p_T > 55$ GeV.

Table 5.1.: Triggers used for the selection of events containing electrons and muons. The p_T thresholds are required at the trigger level, which is determined during operation of the detector.

year	2015	2016
e	$p_T > 24$ GeV and Medium id. level ¹	$p_T > 60$ GeV and Medium id. level ⁴ $p_T > 140$ GeV and Loose id. level ⁵
	$p_T > 60$ GeV and Medium id. level ²	
	$p_T > 120$ GeV and Loose id. level ³	
μ	$p_T > 50$ GeV ⁶	$p_T > 50$ GeV ⁶

Event quality requirements

Events are required to be in a list of good data collection runs, which are determined centrally for all data-taking. Also, the detector is required to be fully operational without errors. Another requirement is that a primary vertex was found in the detector. This means that the reconstruction found an actual physical interaction in the central region surrounded by the detector. Events without a primary vertex might be triggered, but are of no use for the analysis.

Geometrical acceptance requirement

In the transition region between the barrel and the endcaps, which corresponds to values $1.37 < |\eta| < 1.52$, the lack of full instrumentation results in lower electron identification performance. By

¹e24_lhMedium_L1EM20VH

²e60_lhMedium

³e120_lhLoose

⁴e60_lhMedium_nod0

⁵e140_lhLoose_nod0

⁶HLT_mu50

excluding electrons released into that region, this effect can be suppressed. Moreover, all electrons are required to be in the $|\eta| < 2.47$ region, corresponding to an exclusion of electrons in proximity to the beamline, where identification is also not optimal. This is due to the fact that the acceptance of the tracking system ends at $\eta = 2.47$, and electron identification incorporates both tracks and electromagnetic showers in the calorimeter. In the muon channel, no explicit requirements on η are set, the requirement of a *combined* muon (see Section 3.2.4) and the trigger impose an implicit cut of $|\eta| < 2.4$.

Identification and isolation

Using centrally provided prescriptions [31], particles can be identified and the confidence in the identification result can be established. In the case of electrons, a likelihood based approach is leveraged (see Section 3.2.4). For the signal selection, the "LHTight" operating point is chosen. For muons the "High- p_T " working point configuration is used as a criterion. To suppress signatures which are similar to actual electron and muon signatures, but stem from jets, the leptons in the event are required to be isolated (see Section 3.2.4). For electrons the "loose" isolation working point is used, for muons the "loose" working point is used, but isolation is required only for the track.

Explicit track requirements

Both electrons and muons are subjected to a cut on the absolute d_0 significance $|d_0/\sigma_{d_0}|$, which is the magnitude of the impact parameter of the particle track relative to the center of the detector. Using this cut reduces the contribution from reactions which are displaced from the primary vertex. It also suppresses the decay products of long-lived particles, which might be produced and decay only after propagating over a certain distance. The value of this parameter is constrained to $|d_0/\sigma_{d_0}| < 5$ for electrons and $|d_0/\sigma_{d_0}| < 3$ for muons. For muons, an additional cut on $|z_0 \sin \theta| < 0.5$ mm is made. The goal is to suppress cosmic muons, which can pass the detector at virtually any point. To not overly cut on muons in the forward region, where track resolution is lower, the value of z_0 is multiplied by $\sin \theta$. This effectively loosens the cut for muons with θ diverging from $\pi/2$. Also, muons are required to be reconstructed from the combination of a track in the inner detector and a matching track in the muon system.

Additional lepton vetoes

A loosened selection is applied to find additional muons and electrons in the events. For this, the identification requirements for electrons are dropped to "LHMedium", while muons are required to pass either the "High- p_T " or "MediumID" working points. Also, the p_T requirements for both types of particles are dropped to $p_T > 20$ GeV. As a consequence, the signal selection is a subset of

5. Analysis procedure

this loosened selection. If an additional lepton of either type is found, the event is discarded. This procedure suppresses events coming from $t\bar{t}$ and Diboson and neutral current Drell-Yan processes.

Missing transverse energy and transverse mass requirements

As described in Section 2.3, the variable of interest for this search is the transverse mass m_T of the lepton-neutrino pair. The p_T of the neutrino is replaced with the missing transverse momentum E_T^{miss} calculated for the event (see Section 3.2.5).

An additional selection is made based on m_T and E_T^{miss} . The event is required to have $E_T^{\text{miss}} > 65$ GeV for the electron channel, and $E_T^{\text{miss}} > 55$ GeV for the muon channel. The values are chosen to match the respective p_T values mentioned before, which avoids imposing a preference for imbalanced event topologies. Finally, the event is required to have $m_T > 130$ GeV for the electron channel and $m_T > 110$ GeV for the muon channel. The m_T threshold corresponds to $m_T > p_T^l + E_T^{\text{miss}}$. This cut on m_T suppresses on-shell events from the Standard Model Drell-Yan processes, which yield transverse masses below 100 GeV.

The full set of selection criteria is summarized for electrons and muons in Table 5.2.

Table 5.2.: All selection criteria being applied on electrons and muons for the signal selection.

e channel	μ channel
Event level selection	
Event passed one of the triggers in Table 5.1	
Cleaning of event activity	Cleaning of event activity
Lepton level selection	
$ \eta < 2.47$, crack excluded	"combined" muon
LHTight id.	"High- p_T " working point
$p_T > 65$ GeV	$p_T > 55$ GeV
$ d0/\sigma_{d0} < 5$	$ d0/\sigma_{d0} < 3$, $ z_0 \cdot \sin \theta < 0.5$ mm
"Loose" isolation criteria	"Loose" isolation criteria for track only
lepton-neutrino-pair level selection	
Additional e / μ veto: LHMedium, $p_T > 20$ GeV / medium or high- p_T WP, $p_T > 20$ GeV	
$E_T^{\text{miss}} > 65$ GeV, $m_T > 130$ GeV	$E_T^{\text{miss}} > 55$ GeV, $m_T > 110$ GeV

5.2. Data driven estimation of misidentified background events

Aside from the background processes which were already discussed, and which give rise to real leptons and neutrinos or pairs of charged leptons, there is an additional component. The proton-proton collisions produce events which contain large numbers of jets. An example for this is gluon fusion with subsequent splitting into a pair of quarks or gluons. Since gluons make up a large fraction of protons, this reaction occurs with a high rate. Additionally, the Drell-Yan, Diboson and Top background processes can all produce hadronic final states in the form of jets. Jets can produce signatures in the detector, that are not easily distinguishable from lepton signatures originating from actual leptonic decays. With a "fake" lepton signature, the only missing aspect to mirror the signal signature would then be a high value of E_T^{miss} . This can be caused by a mismeasurement of the energy of one of the participating objects, or by a decay product leaving the acceptance of the detector. This leads to an imbalance in transverse momentum, which in turn is the very definition of missing transverse energy. The resulting background is called "Multijet".

This effect is mostly seen in the electron channel, as jet signatures are much more similar to electron signatures than they are to muon signatures. The muon channel is affected as well, albeit to a lesser degree. This analysis focuses on the electron channel, and includes a simplified treatment of the muon channel, to be able to compare the results to published ones. The Multijet background is therefore neglected for the muon channel.

By having passed the selection procedure described in Section 5.1, events originating from this Multijet background are already suppressed, a contribution to the signal region remains in spite of this. In contrast to other types of backgrounds, modeling of the underlying QCD processes is not precise enough to use a Monte Carlo simulation. Therefore, another approach has to be chosen. A procedure called *matrix method* is employed to estimate the contribution of the QCD background in the signal region. It uses the collected data itself, as well as a set of transition probabilities, that can be directly obtained with the help of data,

The starting point is a relationship between the quantities of interest, namely N_R and N_F , corresponding to the number of "real" and "fake" electrons which are contained in the data. "Real", in this case, refers to electrons which are not misidentified, but which do not have to be signal events a priori. "Fake", on the other hand, refers to an electron signature which originates from a jet. To estimate these numbers, two identification levels are defined, the definition of "tight" electrons corresponds to the signal selection, while the definition of "loose" electrons corresponds to the identification level which is established at the trigger level. N_T is the number of electrons accepted in "tight", and N_L is the number of electrons accepted only in "loose". The lower identification level "loose" is exclusive with respect to the "tight" level, which means that N_L counts the number of electrons which are in "loose", but are not contained in "tight". The "real" and "fake" efficiencies

$$\epsilon_F = \frac{N_{\text{tight}}^{\text{fake}}}{N_{\text{loose}}^{\text{fake}}} \quad \epsilon_R = \frac{N_{\text{tight}}^{\text{real}}}{N_{\text{loose}}^{\text{real}}} \quad (5.1)$$

can then be introduced. ϵ_R refers to the probability of "real" electrons passing the "tight" identification

5. Analysis procedure

criteria, while ϵ_F quantifies the probability for "fake" electrons to make it into the same selection.

Using these efficiencies, the migration between the described identification levels can be written as

$$\begin{pmatrix} N_T \\ N_L \end{pmatrix} = \begin{pmatrix} \epsilon_R & \epsilon_F \\ 1 - \epsilon_R & 1 - \epsilon_F \end{pmatrix} \begin{pmatrix} N_R \\ N_F \end{pmatrix}. \quad (5.2)$$

The interesting quantity is $\epsilon_F N_F$, which is the number of "fake" electrons which make it into the "tight" identification level, which is identical to the signal selection. It therefore describes the actual contribution from Multijet events. By inverting Equation (5.2) it follows that

$$\begin{pmatrix} N_R \\ N_F \end{pmatrix} = \frac{1}{\epsilon_R(1 - \epsilon_F) - \epsilon_F(1 - \epsilon_R)} \begin{pmatrix} 1 - \epsilon_F & -\epsilon_F \\ \epsilon_R - 1 & \epsilon_R \end{pmatrix} \begin{pmatrix} N_T \\ N_L \end{pmatrix} \quad (5.3)$$

which gives rise to

$$\epsilon_F N_F = \frac{\epsilon_F}{\epsilon_R - \epsilon_F} (\epsilon_R N_L + N_T(\epsilon_R - 1)) \quad (5.4)$$

Note that now all quantities required to calculate $\epsilon_F N_F$, namely ϵ_F , ϵ_R , N_L and N_T , are measurable.

The first step in producing the estimate is to define the identification levels. For the analysis of electrons, "Tight" is set to the "LHTight" identification criterion with an additional requirement of isolated electrons. The "loose" level is ideally chosen to be as loose as possible, meaning the trigger level. Events with electrons identified at looser levels than required by the trigger are not recorded. In order to preserve statistics, only un-prescaled triggers are eligible. The lowest un-prescaled trigger incorporating the "LHLoose" identification level comes with a cut on the transverse momentum at $p_T > 145$ GeV. Below that, the lowest un-prescaled trigger requires the "LHMedium" identification level. To optimize the procedure, both triggers are combined to make up the "loose" identification level in the sense of the matrix method.

5.2.1. Real efficiency

The real efficiency ϵ_R is determined by selecting events fulfilling the "tight" and "loose" criteria in the Monte Carlo samples, while ensuring the selected electrons are "real" by matching reconstructed objects with the truth level objects, that are contained in the samples. The matching is done by requiring the reconstructed electron to be within a cone of

$$\Delta R = \sqrt{\Delta\phi^2 + \Delta\eta^2} < 0.2. \quad (5.5)$$

This enables a pure real electron selection. The selected electrons are filled into histograms in the properties p_T , η and $\Delta\phi(e, E_T^{\text{miss}})$ for the two identification levels. $\Delta\phi(e, E_T^{\text{miss}})$ is the angle between the electron trajectory and the assumed direction associated with E_T^{miss} in the transverse plane. For events originating from actual decays into an electron-neutrino pair, this separation angle is most likely to be close to π , corresponding to a back-to-back topology.

5.2. Data driven estimation of misidentified background events

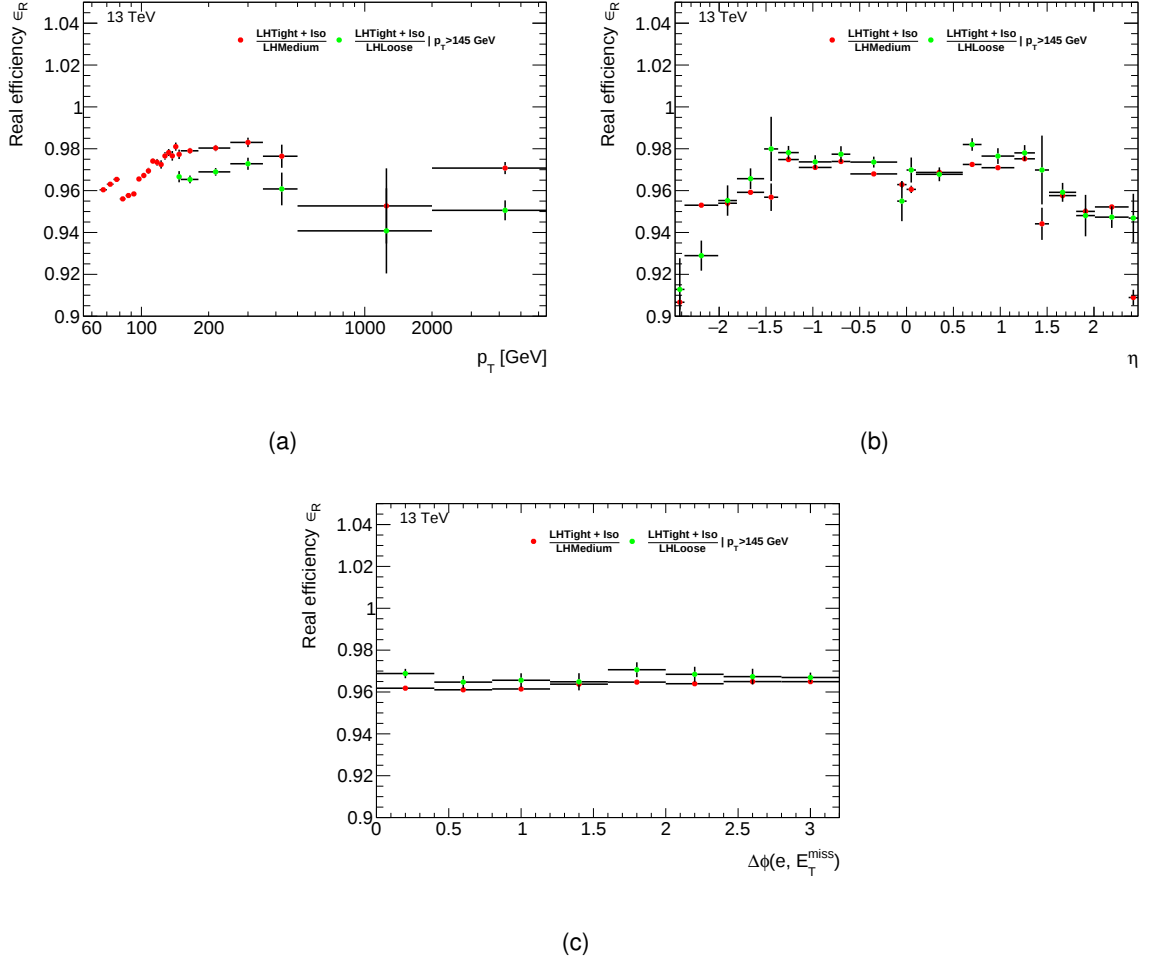


Figure 5.1.: Distributions of the "real" efficiencies, for the "LHLoose" and "LHMedium" lower identification requirement, as a function of various variables.

By dividing the "tight" and "loose" electron histograms for each variable by one another, the real efficiency can be determined with a dependence on those three variables. This is shown in Figure 5.1. The plots show the "real" efficiencies. The combination of "LHMedium" and "LHLoose" is visualized, the efficiency from "LHLoose" is shown for the full p_T spectrum.

Figure 5.2a shows the η distribution of the "real" efficiency. The efficiency is maximal in the central region of the detector, and falls towards the forward regions. The deviation between the two identification levels is less pronounced here, especially for small values of η . Figure 5.2b shows the "real" efficiency as a function of p_T . The efficiency increases up to the threshold of $p_T = 145$ GeV. Above this threshold, a deviation between the two levels of about 2% can be seen. The higher value for the tighter identification level is expected, as the probability for an electron fulfilling "LHMedium" to also fulfill "LHTight" is higher than with "LHLoose" as a starting point. The absolute size of the deviation is relatively small, since the identification criteria are designed to have high signal efficiency. Figure 5.1c shows the "real" efficiency as a function of the angular difference in the transverse plane, which was discussed before. The dependence on this angular difference is limited to the sub-percent

5. Analysis procedure

level, as is the difference between the different identification levels. The "real" efficiency is only and retrieved as a function of p_T and η .

5.2.2. Fake efficiency

A similar approach is used to determine the fake efficiency. Here, the selection is done in a special kinematic region, that is enriched in Multijet events. The goal of this pre-selection is to enhance the number of fake electrons, that goes into the subsequent ratio. The aforementioned cuts on E_T^{miss} or m_T are not applied, as the fake electrons originate mostly from the lower end of the energy spectrum. This, however, opens up the selection to events stemming from on-shell decays of Z^0 and $W^\pm$ bosons, which introduces "real" electrons in the selection. To suppress the Z^0 contribution, the events are scanned for an electron pair, which is required to meet "LHMedium" identification and have at least $p_T > 20$ GeV each. Another requirement is applied by looking for electrons which meet the "LHLoose" criteria, and whose invariant mass is within 20 GeV of the mass of the Z^0 boson at $M_{Z^0} = 91.2$ GeV, and rejecting them. In order to reduce the contribution from $W^\pm$ decays, a veto of $E_T^{\text{miss}} > 60$ GeV is imposed. In order to estimate the amount of "real" electrons which make it into the enriched selection despite the measures mentioned above, the same selection is applied on the Monte Carlo samples. The resulting distributions are subtracted from the ones determined in data prior to the division. The resulting efficiencies as a function of different variables can be seen in Figure 5.2.

The efficiency shown in Figure 5.2a as function of η is relatively flat. The η distribution shows a negative value for the "LHMedium" variant of the "fake" efficiency at the bin at -1.5 , and a strongly decreased value at $+1.5$. This is due to the fact, that these bins cover the transition region between the barrel and the endcaps. In this region, performance identification of the electrons is lowered, and the region is excluded via cuts on η , as described in Section 5.1. The negative value is an artifact of the "real" contribution subtraction. If the "real" estimation from the simulation is above the observed event count in the data, the result of the subtraction can be negative. This can occur due to statistical fluctuations as a result of low numbers of events in these bins. The reason for this is that the cuts on η are applied on the value determined from the second layer of the calorimeter, while the value shown in Figure 5.2a is a combined best estimate. Therefore, only events with a discrepancy between these two η -values migrate to the two bins in question, explaining the low population. Figure 5.2c shows a clear dependence of ϵ_F on $\Delta\phi(e, E_T^{\text{miss}})$ for the "LHMedium" variant.

All distributions in Figure 5.2 show a clear separation between the two lower identification levels, "LHMedium" and "LHLoose". Figure 5.2b shows the "LHMedium" variant for $p_T < 145$ GeV and "LHLoose" for higher values. The "fake" efficiency is at around 0.6 for "LHMedium", and about 0.2 for "LHLoose". ϵ_F is determined dependent on p_T and $\Delta\phi(e, E_T^{\text{miss}})$. The size of the deviation between the two levels is significantly larger as for the "real" efficiency. This is expected, since for "fake" electrons, the increase in identification strictness should result in higher rejection. The "fake" efficiency is stored and retrieved as a function of $\Delta\phi(e, E_T^{\text{miss}})$ and p_T .

5.2. Data driven estimation of misidentified background events

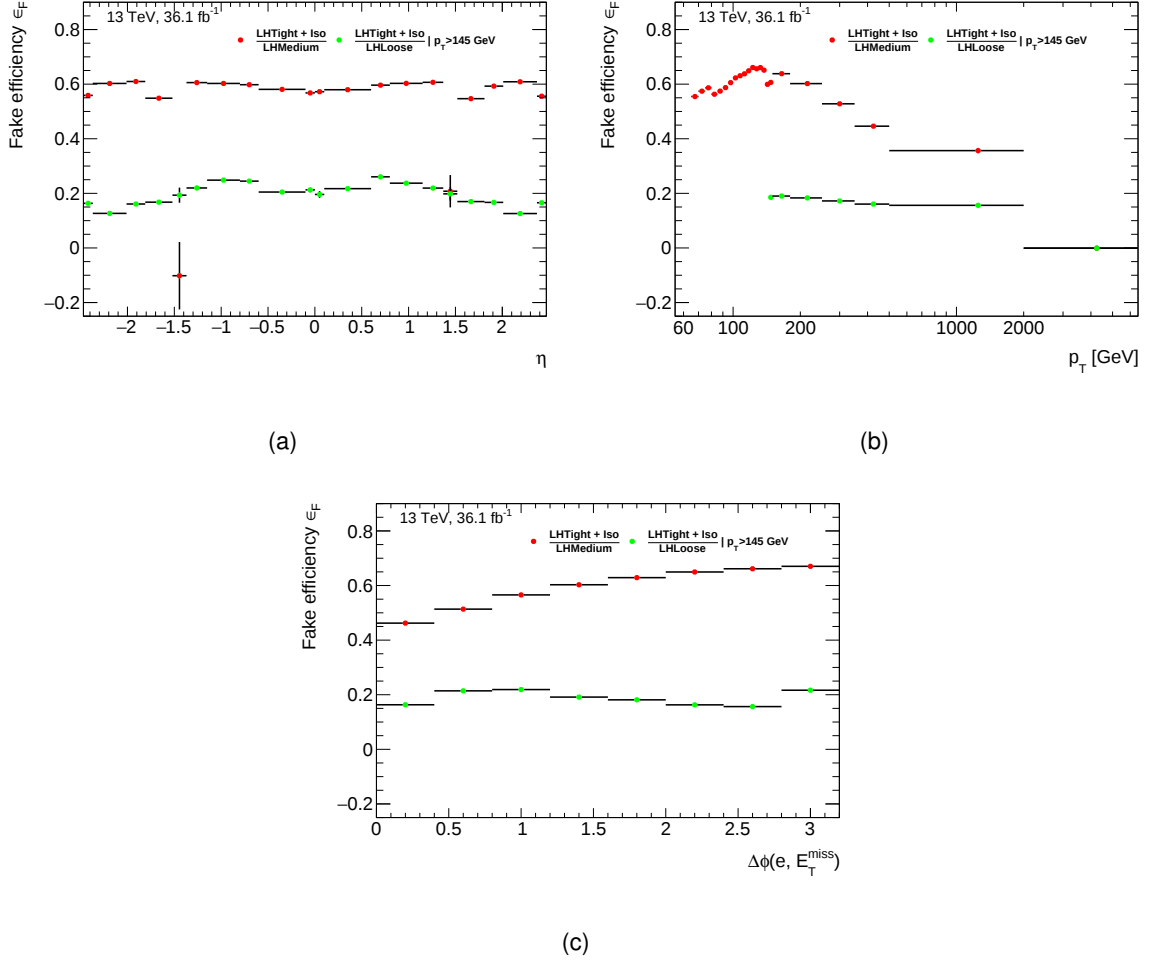


Figure 5.2.: "Fake" efficiency distribution as a function of various variables. The efficiencies are used using the full integrated luminosity at $\sqrt{s} = 13 \text{ TeV}$, 36.1 fb^{-1} .

Having determined the real and fake efficiencies, the last step to produce an estimate for the QCD background is to implement the two identification levels in the signal region, where the contribution from "fake" electrons is supposed to be studied. Using the same "LHTight" and isolated and "LHLoose" / "LHMedium" criteria, each selected electron can then be assigned a value for ϵ_F and ϵ_R , depending on the kinematic variables. A weight based on Equation (5.4) can then be determined, depending on whether the electron is marked as "tight" or "loose".

$$w_L = \frac{\epsilon_F \epsilon_R}{\epsilon_R - \epsilon_F} \quad w_T = \frac{\epsilon_F (\epsilon_R - 1)}{\epsilon_R - \epsilon_F} \quad (5.6)$$

This weight is used when filling the various properties of the electron into the histograms that are used in subsequent analysis steps. The resulting reweighted data sample corresponds to the estimation of the contribution of "fake" electrons in the signal region, and is subsequently treated as the background sample "QCD".

5.3. Kinematic distributions

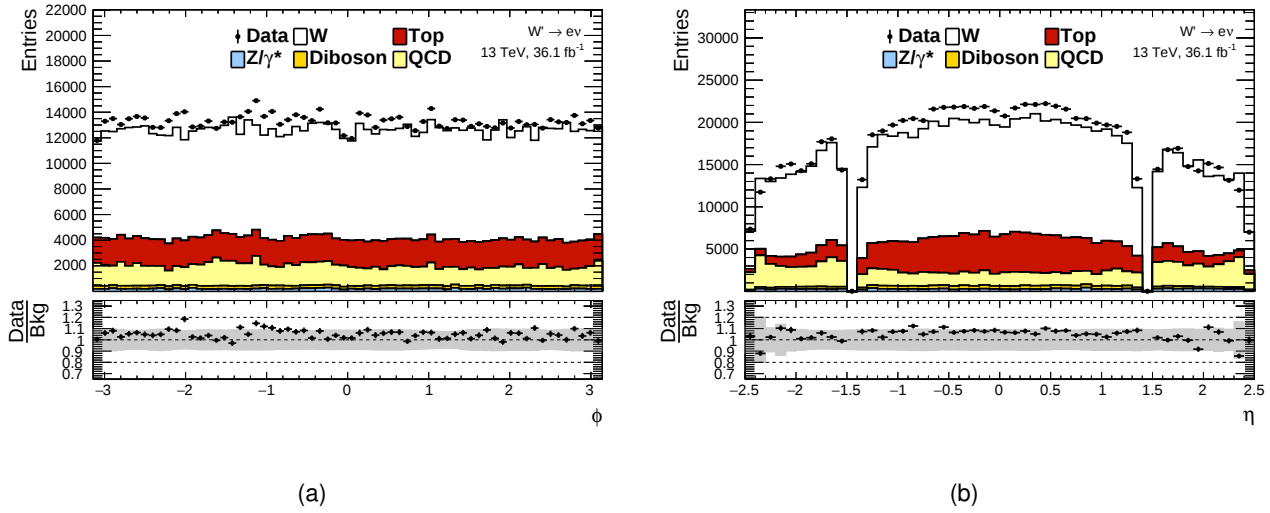


Figure 5.3.: Electron angular distributions in ϕ and η for the signal region.

Using the selection procedure described in Section 5.1, histograms of all interesting properties can be derived using the various background Monte Carlo samples, as well as the sample used for the Multijet background. For improved visualization the different background samples are combined into their source groups, as defined in Table 4.3.

To gauge the modeling of the collected data by the background expectations, these kinematic distributions are a useful tool. The kinematic parameters of the leptons that are shown are the ϕ angle and the pseudorapidity η , which describe the transverse and longitudinal direction of the particle, respectively. Also shown is the transverse momentum p_T of the leptons. Additionally, the value for the missing transverse momentum in the event is shown. The figures show the respective variable and the corresponding number of events for every interval of the histogram.

5.3.1. Electron channel

The contribution from different background sources to the electron channel signal selection can be seen nicely by looking at the ϕ distribution in Figure 5.3a, which contains the information from all energetic configurations.

The ϕ distribution is flat across its value range, which is expected since no preferred decay direction exists in the transverse plane, corresponding to the isotropy of space, and no construction features of the detector introduce any structures in this variable. The distribution of η in Figure 5.3b shows a slight decrease towards higher absolute values. The distribution also features two regions of negligible entries, around $1.37 < |\eta| < 1.52$, which corresponds to the so-called "crack" between the central region and endcap region of the detector. Since identification efficiency is lowered in these transition

regions, they are excluded from the signal region, explaining the structures.

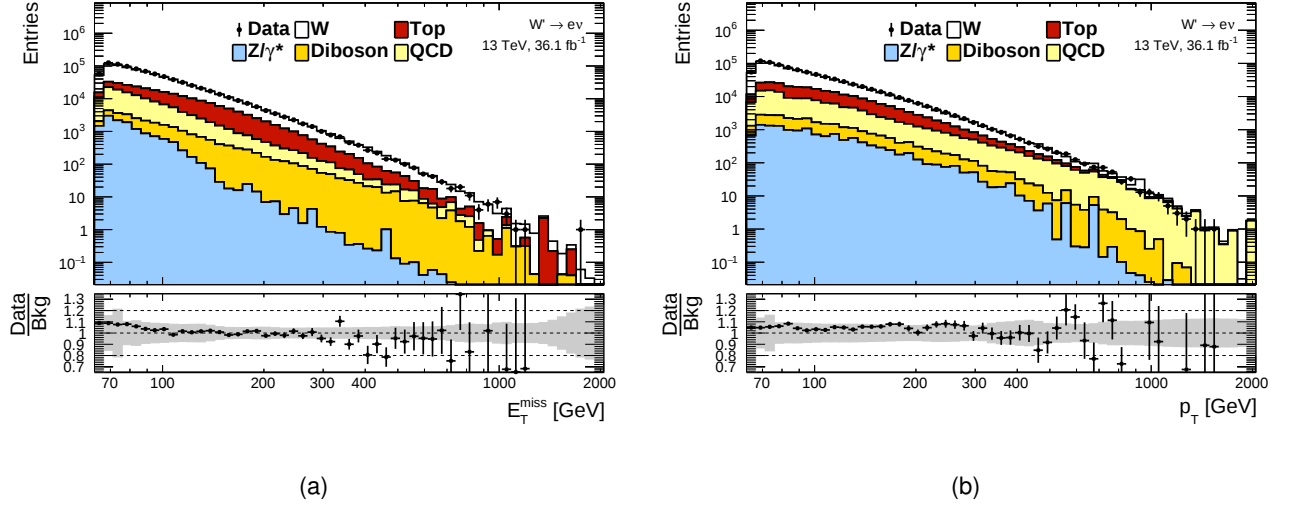


Figure 5.4.: Electron p_T and E_T^{miss} distributions in the signal region.

The transverse momentum as well as the missing transverse energy show steeply falling spectra, shown in the double logarithmic plots in Figure 5.4a and Figure 5.4b.

Along with the background sources, the data is shown as dots with their statistical errors drawn as bars. The lower part of the plots shows the ratio of the number of observed events in the data, and the sum of all backgrounds. Aside from that, a gray band indicates the squared sum of all systematic uncertainties considered for the background samples, which will be discussed in Section 5.4. The ratio is flat for ϕ , η and p_T , the latter distributions runs out of statistics starting at around 500 GeV. The ratio in ϕ lies relatively close to unity, leading to the conclusion that when looking at the entire range of energetic configurations, the data is described well by the background estimations. There is, however, a problematic structure at small values of E_T^{miss} , where the ratio rises as far as 1.1. This effect is not fully understood, but is thought to be caused by a mismodeling of the impact parameter in some of the simulated samples and resolution effects in the calculation of the soft term of the E_T^{miss} . Nevertheless, the systematic uncertainties cover the deviation of the ratio from unity, meaning that proper treatment of the systematic uncertainties in later steps of the analysis can account for and control the effect seen at low E_T^{miss} .

5.3.2. Muon channel

Analogously to the electron channel, the same kinematic distributions can be studied for the muon channel. Figures 5.5 and 5.6 show the distributions for ϕ , η , E_T^{miss} and p_T for the muons. In contrast to the electron channel, the ϕ distribution has a structure imposed on the otherwise flat shape. This is due to the additional material of the magnets and the support structure of the detector, which the muons have to cross on their path to the muon system. This additional material affects the reconstruction efficiency, and reduces the event yield in these regions.

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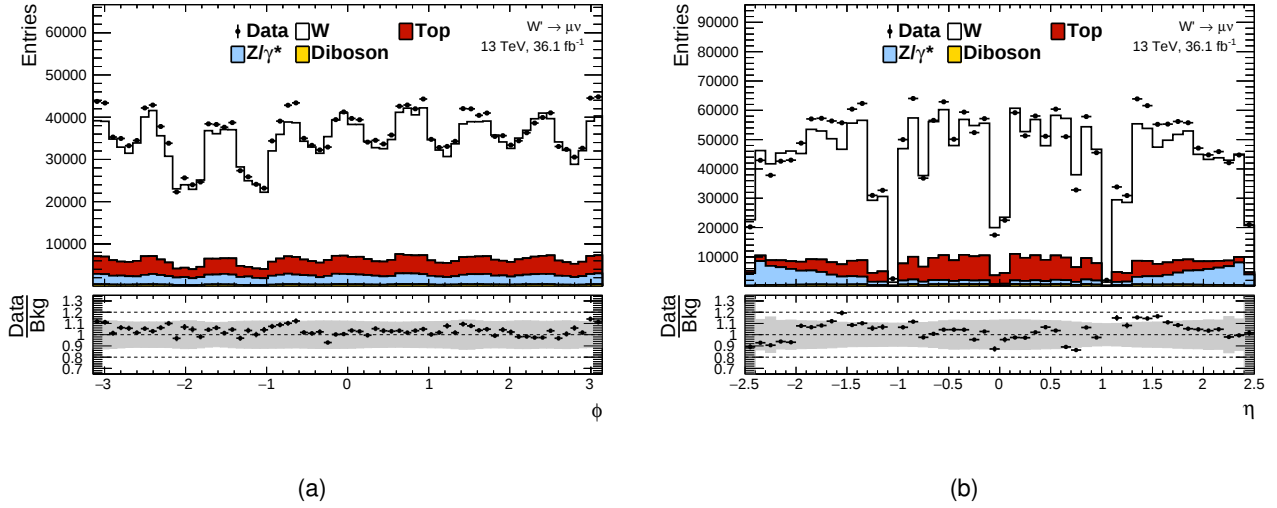


Figure 5.5.: Muon angular distributions in ϕ and η for the signal region.

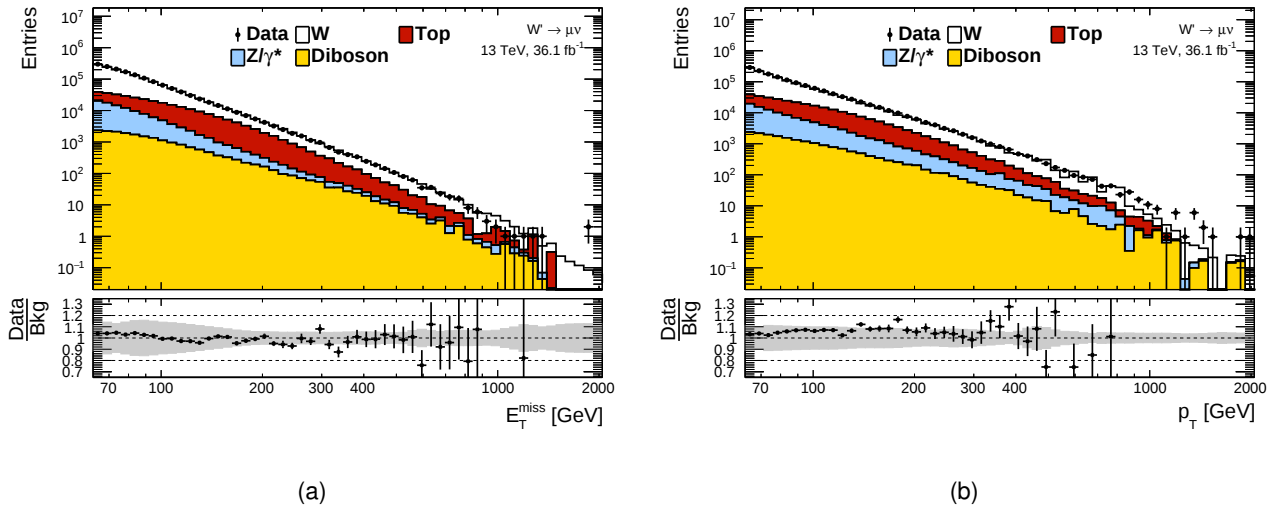


Figure 5.6.: Muon kinematic distributions in E_T^{miss} and p_T in the signal region.

The η distribution features the same drops in the transition region between the central and the endcap region at around $|\eta| \approx 1.5$ as seen in the electron channel, while also including a drop at $\eta \approx 0$, where the muon system has a gap.

The E_T^{miss} and p_T distributions show similar falling spectra as seen in the electron channel. The agreement between data and background seems marginally better than in the electron channel. The E_T^{miss} distribution does not show a deviation at low values. The muon channel is, however, lacking any sort of scale factor treatment of the Monte Carlo samples, which means that the compatibility of the background and data might change if those were introduced.

5.4. Treatment of systematic uncertainties

Aside from statistical uncertainties, which are an intrinsic property of every type of measurement, systematic uncertainties also need to be taken into account. These uncertainties need to be determined externally, since they can not be directly inferred from data. Systematic uncertainties can be structured into two categories, the first one groups uncertainties coming from a lack of knowledge from a theoretical perspective. The second category originates from the reconstruction and analysis procedure, meaning that it is connected to the experiment itself.

Systematic uncertainty sources might impact the description of the background, as well as the signal expectation, both of which are derived using Monte Carlo simulations. If a source affects both the signal and the background expectation, it needs to be applied in a correlated way. For this analysis, only the impact on the background expectation is taken into account, variation of the signal expectation is neglected. For the centrally developed software packages, systematic uncertainties are provided where applicable.

5.4.1. Theoretical uncertainties

The Monte Carlo simulation procedures use parton distribution functions to generate events in proton-proton collisions, which cannot be calculated without experimental input. The uncertainties on the specific value of the PDF needs to be accounted for, as well as the concrete choice of PDF.

The k -factor corrections, which are applied to the Monte Carlo simulation in order to increase their precision, include several uncertainties related to the perturbation theory treatment used in the determination of the mass dependent scaling factors. Since these factors are used to estimate the background, they need to be accounted for. Also, cross-sections being used in the generation of physical processes are susceptible to systematic shifts, which are considered as well.

5.4.2. Experimental uncertainties

The second category are uncertainties which are picked up during reconstruction of detector signals or the analysis itself. The energy scale and resolution corrections, which are applied to electrons in data and simulation, respectively, introduce uncertainties. The same applies to jet energy scale and resolution, which are part of the calculation of E_T^{miss} . Additionally, the efficiencies which are applied for the reconstruction, trigger, isolation and identification steps, also introduce systematic uncertainties. They can be estimated for each of the efficiencies separately. Aside from the jets, other effects influence the calculation of E_T^{miss} , specifically the inclusion of soft-objects, meaning the sum of objects not associated with the primary physics objects of the selection.

5. Analysis procedure

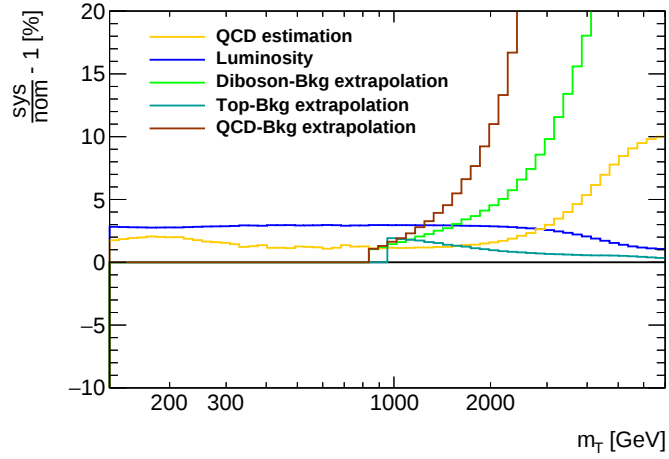
Due to the extrapolation procedure employed to improve the samples for the Multijet, Top and Diboson backgrounds, another systematic uncertainty will be introduced in Section 5.5. As described in Section 3.2.6, the integrated luminosity $\mathcal{L}$ that the collected data sample amounts to, is a quantity which needs to be measured. The Monte Carlo samples are scaled to the measured integrated luminosity, taking into account the respective cross-sections. Thus, the uncertainty on the value of the integrated luminosity is considered. For this, the impact of a variation of the scale of the Monte Carlo samples, amounting to 3.2 %, is studied.

The Multijet background is not directly affected by the luminosity uncertainty, since it uses the data sample itself for reweighting, which is at the correct integrated luminosity. On the other hand, the Multijet background also has sensitivity towards the measured integrated luminosity indirectly, due to the subtraction of the "real" contribution in the Multijet enriched region. This effect is accounted for, by applying a constant uncertainty of 15 % on the Multijet background alone. Finally, the procedure to account for pile-up effects, as discussed in Section 4.4.3, also includes uncertainties on the data scale factor.

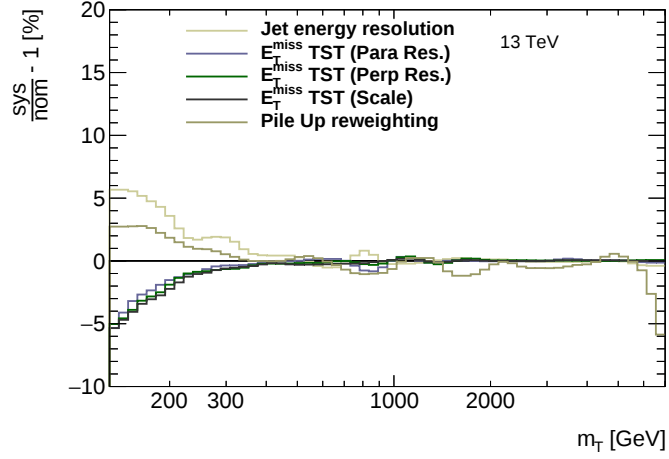
All uncertainties are evaluated regarding their impact on the total sum of the background, meaning that a large variation in one of the smaller backgrounds might still only have a negligible impact. Only systematic uncertainties whose impact on the total background is in excess of 3 % are considered.

5.4.3. Systematic uncertainties as a function of transverse mass

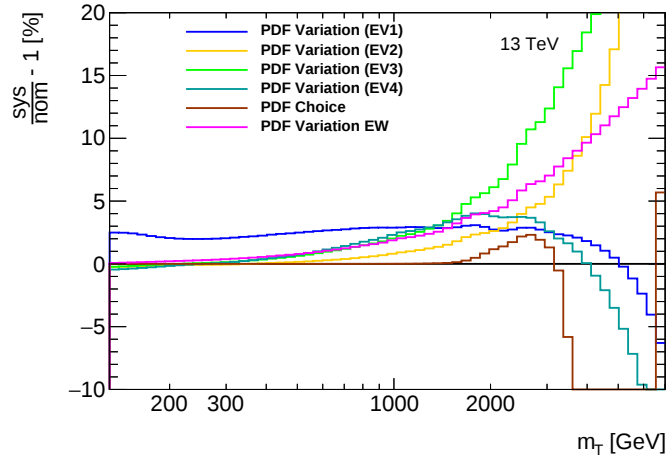
Figure 5.7 shows the various systematic uncertainties on the background that are accounted for. Figure 5.7b shows the impact of the electron and jet energy scale and resolution, respectively, as well as three E_T^{miss} related uncertainties. The electron energy scale increases with m_T and reaches about 8.5 %, before dropping at the end of the spectrum. The jet energy scale, as well as the E_T^{miss} systematic uncertainties have effects on the 5 % level for low m_T , dropping to sub-percent levels for higher masses. Figure 5.7a shows the luminosity uncertainty, the flat uncertainty assumed for the Multijet background, as well as the uncertainties derived for the extrapolation procedure of the Multijet, Diboson and Top backgrounds, described in Section 5.5. The latter three start at around $m_T = 1$ TeV, since the extrapolated distribution replaces the Monte Carlo prediction above a certain value of m_T .



(a) Luminosity related and extrapolation uncertainties.



(b) Energy scale and resolution and Pile-up uncertainties.



(c) Theoretical uncertainties.

Figure 5.7.: Plots of the theory, luminosity, extrapolation and energy related systematic uncertainties which are considered for the analysis for the electron channel. The plots show the relative size of the deviation from the nominal background estimate in percent for different values of m_T .

5. Analysis procedure

Figure 5.7c shows the theoretical uncertainties. The variation of the PDF value comes in the form of a set of uncorrelated eigenvectors, four of which are kept for the analysis based on their impact, and shown here. The uncertainty from the choice of PDF is shown as well, along with the variation of the electroweak correction. All of the aforementioned uncertainties remain below 5 % until the high mass end of the spectrum, where they reach considerable magnitudes.

The uncertainties are determined for all kinematic properties, but since only the m_T distribution is relevant for the search procedure itself, the contributions are not shown for other variables. The systematic uncertainties are considered to be symmetric in the statistical interpretation. The quadratic sum of the systematic uncertainties, however, is shown in Section 5.3 as the up and down variation.

Corresponding plots for the muon channel can be found in Appendix A.3. The theory and E_T^{miss} related uncertainties are considered for the muon channel, as is the pile-up uncertainty. The luminosity uncertainty is assumed to be the same magnitude as in the electron channel. In the electron channel, the Multijet background is excluded from the impact of the variation, since this background is automatically scaled to the correct integrated luminosity. In the muon channel, the Multijet background is not considered, hence the luminosity uncertainty results in a flat 3.2 % variation across the entire mass range. The aforementioned uncertainties are assumed to be fully correlated to the electron channel.

5.5. Background extrapolation

Since the combined background samples for the Diboson and Top backgrounds, as well as for the Multijet background have drastically falling statistics towards higher values of m_T , a fit based approach is taken to improve the distributions there. The extrapolation procedure is based on software not developed for this thesis [54, 53].

For this, two fit functions are defined as

$$\begin{aligned}\frac{dN}{dm_T} &= c_1 \cdot m_T^{c_2 + c_3 \log m_T} \\ \frac{dN}{dm_T} &= \frac{c_4}{(m_T + c_5)^2}\end{aligned}\tag{5.7}$$

where c_1 to c_5 are free parameters. Both functions are not physically motivated, instead they are chosen to approximate the expected distributions reasonably. The functions are fitted to a range in the m_T background distribution. The position and size of the fitting range is varied. This step results in a set of fits which describe the underlying distribution to a varying degree. The ensemble of fit functions can then be used to estimate an error on the extrapolation.

The nominal value, meaning the one which is used for the determination of the bin contents in the high mass end of the spectrum, is chosen by one of two methods. For the Diboson and Top backgrounds, the χ^2/N_{dof} is calculated, and the fit with a value closest to unity is chosen as the central value. Subsequently the uncertainty is determined on a bin by bin basis. For each bin the deviation of all fit functions from the previously chosen central value is calculated, and the one featuring the largest deviation is taken as the uncertainty for that bin. In effect, this takes the envelope of the ensemble of fitted functions as the uncertainty.

This procedure was found to be unstable with respect to the choice of starting parameters for the fitting window position and size, which is undesirable. To improve this, an alternative procedure [53] was implemented. This alternative approach modifies the selection of the central value and the selection of fitted functions being considered for the envelope by using the quantity

$$w = \frac{1}{n_b} \sum_{b > b_s}^{N_b^{\text{est}} > 0} \frac{(N_b^{\text{est}} - N_b^{\text{fit}})^2}{\sigma_b^2} \quad (5.8)$$

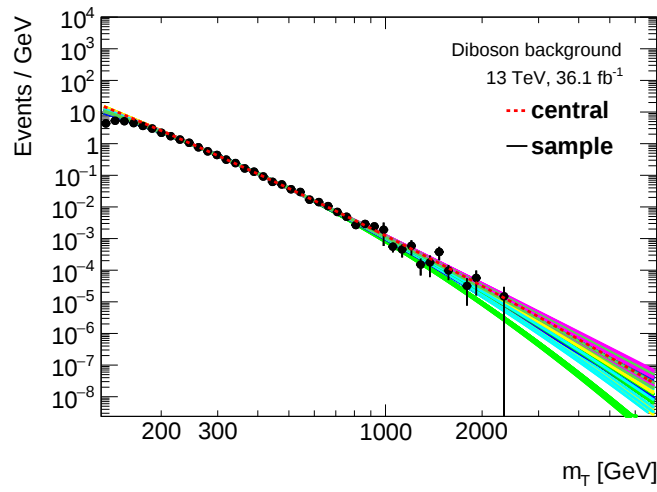
where n_b is the number of bins, N_b^{est} and N_b^{fit} are the bin content of the Monte Carlo histogram and the fitted function value at the center, respectively, and σ_b is the statistical uncertainty of the MC estimate. The fitted function with the lowest value of w is then chosen, while all fitted functions with $w < 1.5$ are considered for the envelope, which is then determined as described before.

The result of the fitting with varying window position and size can be seen in Figure 5.8. The markers and their error bars correspond to the Monte Carlo estimate derived directly from the sample, while the dashed red line indicates the fitted function which was chosen as the central value. Around the central value function the remaining ensemble of fitted functions can be seen in different colors.

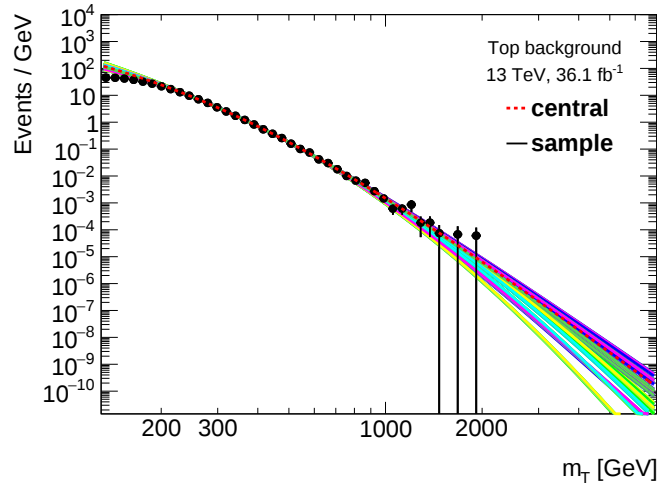
On first look, the extrapolation describes the Monte Carlo prediction reasonably well. However, since the high end of the spectrum has very few events, an evaluation is difficult. The cumulative bin content starting from the high end of the spectrum is plotted in Figure 5.9. The solid markers again show the estimate from the Monte Carlo sample, while the hollow markers correspond to the central value chosen from the ensemble. The red bands indicate the uncertainty derived from the envelope of the fitted functions.

As is visible in Figure 5.9c the central value nicely reproduces the shape of the MC estimate, and continues smoothly towards higher masses. The MC estimate lies within the red uncertainty. Nevertheless, it is sensible to only use the extrapolation for the high mass region it is needed for, relying on the original MC estimate for lower masses. This is achieved by creating a new histogram for the distribution, and selecting a bin, above which the bin content is taken from the extrapolation function. The output of this operation is used in later stages of the analysis.

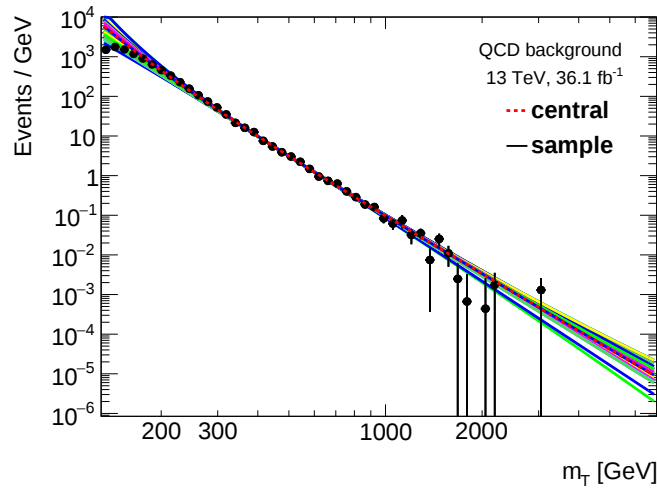
5. Analysis procedure



(a) Diboson background

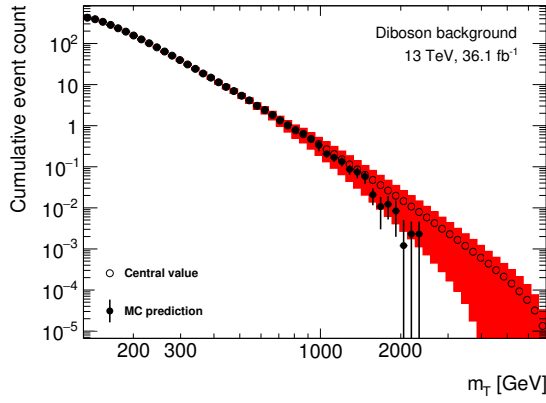


(b) Top background

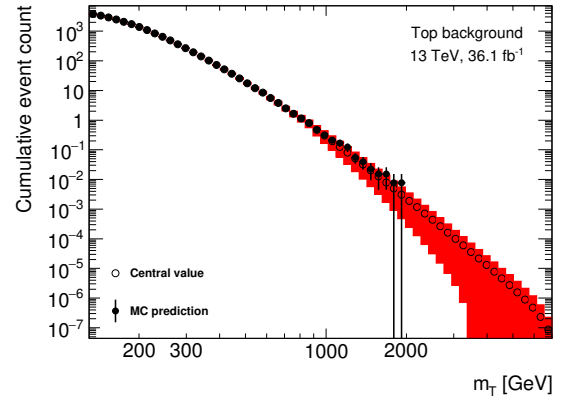


(c) QCD background

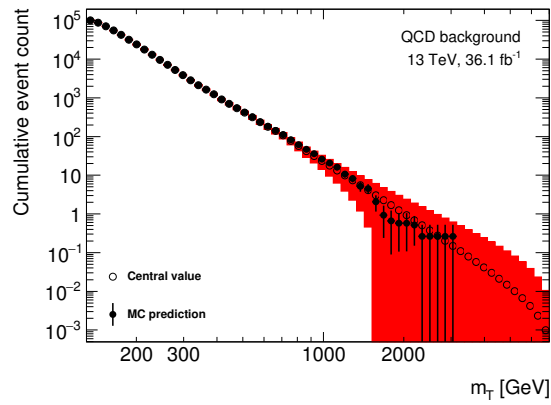
Figure 5.8.: Visualization of the extrapolated backgrounds. The various fit function variations are shown, as well as the central value.



(a) Diboson background



(b) Top background



(c) QCD background

Figure 5.9.: Cumulative histograms, showing the central value of the extrapolation and the determined error, as well as the Monte Carlo prediction.

6. Statistical analysis

After the selection of signal events in the collected data set, and construction of the background expectation in the signal selection, which was described in Chapter 5, the remaining task is the interpretation. This analysis is aimed at finding resonance shaped signals in the transverse mass distribution. The distribution of m_T in data can be seen in Figure 6.1 and Figure 6.2 for the electron and muon channels, respectively.

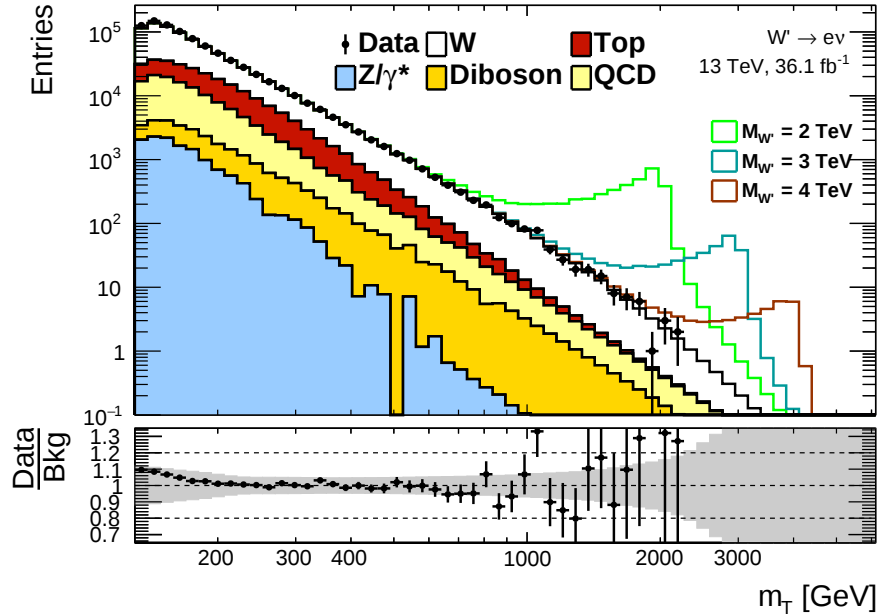


Figure 6.1.: Plot of the m_T distribution in the $W' \rightarrow e\nu_e$ signal selection. The plot shows the data as well as the different background distributions in the upper panel. Also, three signal samples with different pole masses are shown. The lower panel shows the ratio between the data and the sum of backgrounds, as well as the quadratic sum of the considered systematic uncertainties in the band.

The transverse mass distribution for the electron channel shows a steeply falling spectrum towards higher values of m_T . The data distribution is shown as dots, along with the statistical uncertainty. The statistical uncertainties begin to severely impact the distribution at around $m_T \approx 1$ TeV. The different background sources can be seen in varying colors. The dominant background source can be identified as the charged current Drell-Yan process, followed by the Top background. In the electron channel, Multijet background constitutes the third largest contribution, followed by the Diboson and neutral current Drell-Yan processes. The muon channel lacks an estimation of the Multijet background, since

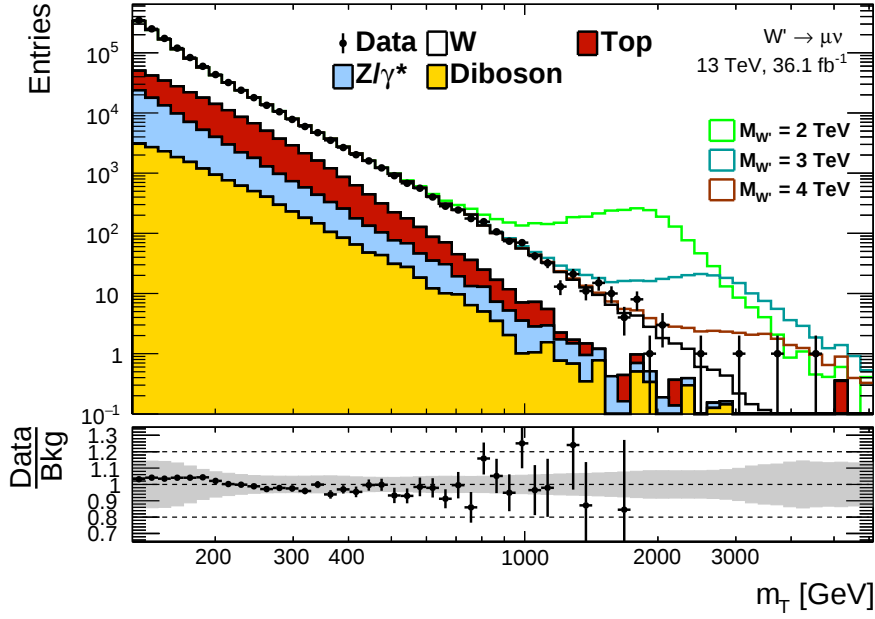


Figure 6.2.: Plot of the m_T distribution in the $W' \rightarrow \mu\nu_\mu$ signal selection. The plot shows the data as well as the different background distributions in the upper panel.

this analysis focuses on the electron channel. The Diboson background has the smallest contribution in the muon channel, whereas the charged current Drell-Yan background appears larger, when compared to the electron channel. This is due to differences in the reconstruction of both types of particles. An event which stems from a dilepton decay only enters the signal region if it has high missing transverse energy. This only occurs, if one of the two produced leptons is miss measured, or escapes the acceptance of the detector. In the case of an electron, reconstruction is done by matching a charged track in the tracking system with a shower in the calorimeter. If the particle track is not covered by the tracking detectors, the calorimeter, which features higher solid angle acceptance, will most likely still register the shower which the electron causes. Without a matching particle track, this object will be reconstructed as a jet, and thus still not increase the value of E_T^{miss} . Muons, on the other hand, are reconstructed by matching a track from the inner detector with one in the muon system, the muon does not create a shower in the calorimeter. If the muon track is not registered, the muon will most likely not be reconstructed at all, thus contributing towards E_T^{miss} .

Before the data set begins to run out of statistics, the ratio of data over background in the electron channel is compatible with unity. This points towards a good description of the data by the background estimate. Only at values $m_T < 200$ GeV does the ratio deviate from unity, reaching a factor of about 1.1 for the smallest values of m_T . The deviation is directly associated with a similar feature, which is observed in the E_T^{miss} distribution shown in Figure 5.4a. This effect, as was previously mentioned, is thought to be related to a problem in the modeling of the impact parameter in the charged current Drell-Yan background samples, and to resolution effects in the calculation of E_T^{miss} . The band in the lower panel, which shows the quadratic sum of the systematic uncertainties discussed in Section 5.4, covers the difference of the ratio to unity.

The transverse mass distribution for the muon channel has a very similar shape as seen for the electron channel. The muon channel is lacking an extrapolation treatment, which explains the fluctuations in the backgrounds for higher masses, which are not visible in the electron channel. A recent ATLAS publication [50] contains muon distributions featuring scale factors, extrapolations, and an estimate of the Multijet background. The ratio of data over background is slightly worse in spite of the scale factors. The Multijet background appears to be the smallest background source in the muon channel. When comparing the systematic uncertainties of the electron and muon channel, the lower uncertainty at high masses for muons is notable. This is due to the lack of extrapolation uncertainties in the muon channel. Aside from that, the electron channel includes a constant 15 % uncertainty, which is applied on the Multijet background and evaluated for the total background. Since there is no Multijet background estimate for the muon channel, this uncertainty is also not included. As in the electron channel, the systematic band covers the deviations of the data over background ratio to a reasonable degree.

In addition to the data and background distributions, three signal hypotheses for the SSM W' without destructive interference are shown for pole masses. The electron signal shapes look significantly narrower than the ones shown for the muon channel. This is due to the lower energy resolution of the detector for the measurement of muons.

By looking at the distributions and the ratio between data and background, it becomes apparent, that no significant excess is visible in the spectrum. Statistical methods can be used in order to quantify and evaluate the compatibility of the data with the background-only hypothesis.

6.1. Reweighting of signal templates

An expectation of the shape of the signal process in the variable under study is a crucial part for the statistical analysis of the data. For this search the SSM is used as a benchmark, which has a generic Jacobian-Peak signal shape in m_T . To produce signal expectations for a variety of pole masses a reweighting technique is used. For this, a Monte Carlo sample is produced, as described in Section 4.3, with a flat distribution of the transverse mass. The following sections describe how a weight based on the matrix element (see Section 2.3.1) can be used to produce signal expectations for arbitrary pole masses. The weights were developed by Mihail Chizhov [55].

6.1.1. Non-interference reweighting procedure

As discussed in Section 2.3.1, making the assumption that the W' is *right-handed*, the mixed-term can be neglected. The consequence of this is that

$$|\mathcal{M}|^2 - |\mathcal{M}_{\text{SM}}|^2 = |\mathcal{M}_{\text{SSM}}|^2 \geq 0 \quad (6.1)$$

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can not be negative. Equation (6.1) describes the effective contribution of the hypothetical signal process to the amplitude of the matrix element. Therefore, it is clear that the introduction of a SSM W' results only in an increase in the number of events observed, when compared to the background expectation. In order to be capable of comparing the measured data to the background-only and signal-plus-background hypotheses, the deviation of a possible signal signature from the background is required. A histogram with bin contents set to the deviation from the background-only hypothesis is created. Such a histogram is referred to as a signal "template". To obtain a template like this, the flat mass sample can be used. By taking the ratio between the cross-section of this flat distribution and the desired resonance shape, a weight [55] can be introduced which, when applied to the events in the flat sample, will result in a template shape.

$$w(m_{l\nu}) = \frac{\sigma_{W'}(m_{l\nu})}{\sigma_{\text{flat}}(m_{l\nu})} \quad (6.2)$$

This weight is a function of the invariant mass of the lepton-neutrino system $m_{l\nu}$ and depends on the W' pole mass $M_{W'}$. It can be written as

$$w(m_{l\nu}) = \kappa^2 \cdot \begin{cases} \frac{10^{12}}{102.77} \exp\left(\frac{-11.5m_{l\nu}}{\sqrt{s}}\right) \cdot W_{BW} & m_{l\nu} < 299 \text{ GeV} \\ 10^{12} \exp\left(\frac{-16.1m_{l\nu}}{\sqrt{s}}\right) \cdot \left(\frac{m_{l\nu}}{\sqrt{s}}\right)^{1.2} \cdot W_{BW} & 299 \text{ GeV} \leq m_{l\nu} < 3003 \text{ GeV} \\ \frac{10^{16}}{1.8675} \exp\left(\frac{-31.7m_{l\nu}}{\sqrt{s}}\right) \cdot \left(\frac{m_{l\nu}}{\sqrt{s}}\right)^{4.6} \cdot W_{BW} & m_{l\nu} \geq 3003 \text{ GeV} \end{cases} \quad (6.3)$$

where

$$\kappa = \left(\frac{g_{W'}}{g_W}\right)^2 \quad (6.4)$$

is the quadratic ratio of the coupling constant of the W' and the W . $\kappa = 1$ corresponds to equal coupling strength, which is postulated in the SSM. $\sqrt{s}$ is the center-of-mass energy of the pp -system and

$$W_{BW} = \frac{1}{(m_{l\nu}^2 - M_{W'}^2)^2 + m_{l\nu}^2 \Gamma_{\text{rel}}^2} \quad (6.5)$$

describes a Breit-Wigner shape. For this, the relative decay width of the W' ,

$$\Gamma_{\text{rel}} = \frac{3 + (1 + r/2)(1 - r)^2}{4 \times 0.2312 \times (127.918 + 1.45 \log(M_{Z^0}/M_{W'}))} \times \kappa. \quad (6.6)$$

is needed, which also directly depends on κ . Here $M_{Z^0} = 91.2 \text{ GeV}$ is the mass of the SM Z^0 and the parameter

$$r = \min \left\{ \left(\frac{m_t}{M_{W'}} \right)^2, 1 \right\} \quad (6.7)$$

is the inverse squared fraction of mass relative to the top mass, capped at one.

The weight is calculated for every event in the flat sample, by using the truth value for the invariant mass $m_{l\nu}$. The resulting distributions for a signal with specific pole masses are shown in Figure 6.3 for various values of κ . κ depends on the relative coupling strengths of the hypothetical particle under study, and influences its cross-section as well as the width of the resonance. This connection becomes apparent when looking at the evolution of the signal shape depending on the values of κ , shown as different colors from blue to red. Higher values of κ broaden the width of the resonance, while also increasing the number of events across the entire m_T range.

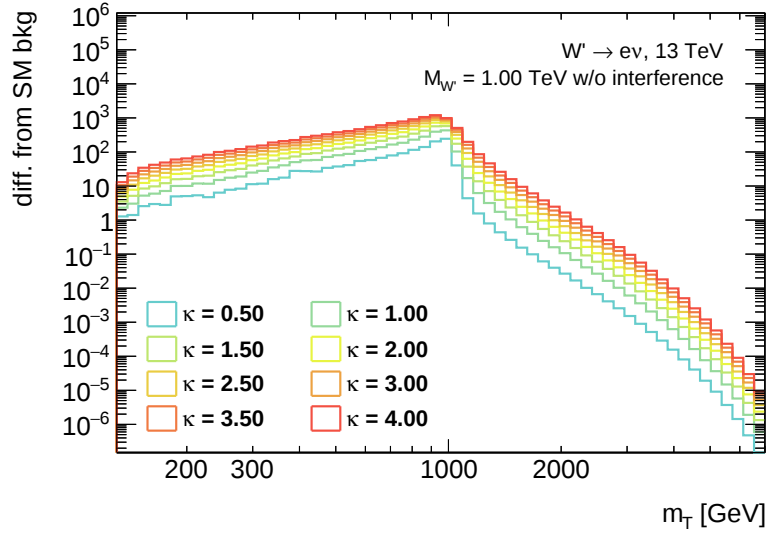


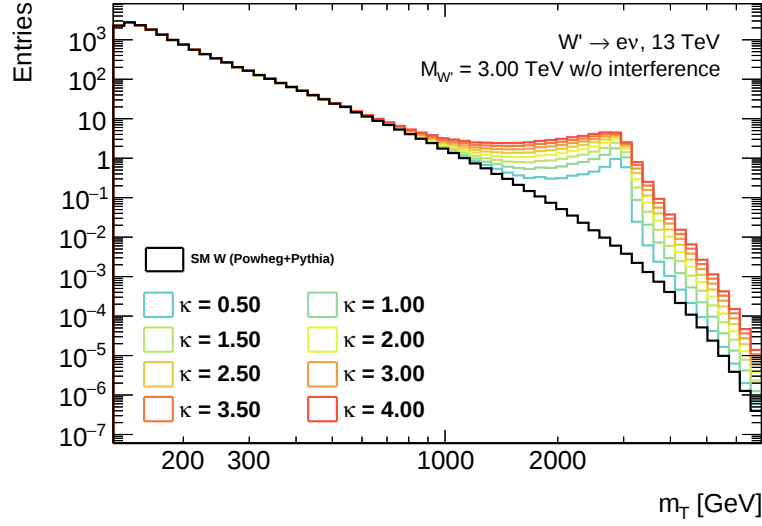
Figure 6.3.: Transverse mass distribution in the electron for a W' pole mass of $M_{W'} = 1$ TeV, excluding the SM contribution and interference effects, as a function of different values of κ . The distributions are scaled to an arbitrary integrated luminosity of 1 fb^{-1} .

To make the interpretation of the actual distributions easier, the W Monte Carlo samples, that are simulated using POWHEG+PYTHIA8 as described in Chapter 4, can be added to the template. This is shown in Figure 6.4. Unfortunately, however, this introduces the problem that the W background sample is simulated at NLO, while the signal sample is simulated at LO. Even though both samples are corrected to NNLO in QCD, only the W background sample also receives a correction to NLO for electroweak contributions (excluding final state radiation), which the signal sample does not (see Chapter 4). Conversely, Figure 6.5 shows the same distributions for a fixed value of $\kappa = 1.0$ and pole masses $M_{W'} = \{1, 2, 3, 4, 5, 6\}$ TeV. Towards higher pole masses, the resonance gives less signal events while the width shrinks on the logarithmic scale.

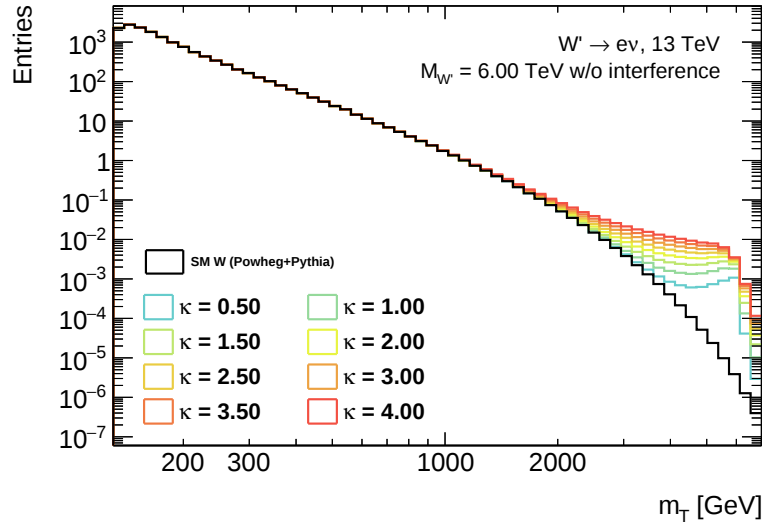
Corresponding distributions for the signal templates determined in the muon channel can be found in Appendix A.1. The plot showing only the deviation of the signal from the background can be found in Figure A.1. When comparing with Figure 6.3, the lower momentum resolution of the muons becomes apparent through the wider resonance shape.

6.1.2. Incorporating interference effects

Making the diametrical choice and assuming a *left-handed* W' , interference between the additional boson and the W boson can occur. Depending on the sign of the coupling constant, this interference effect can be constructive, leading to an increase in the observed number of events. The other possibility is destructive interference, which means that the inclusion of the new particle can reduce the number of observed events, when compared to the background only expectation, and is the one being considered in the following.



(a) $M_{W'} = 3$ TeV



(b) $M_{W'} = 6$ TeV

Figure 6.4.: Transverse mass distributions in the electron channel for different W' pole masses, added to the SM W contribution (POWHEG+PYTHIA8) but excluding interference. The distributions show the results of different values of κ , indicated by the colors. The distributions are scaled to an arbitrary integrated luminosity of 1 fb^{-1} .

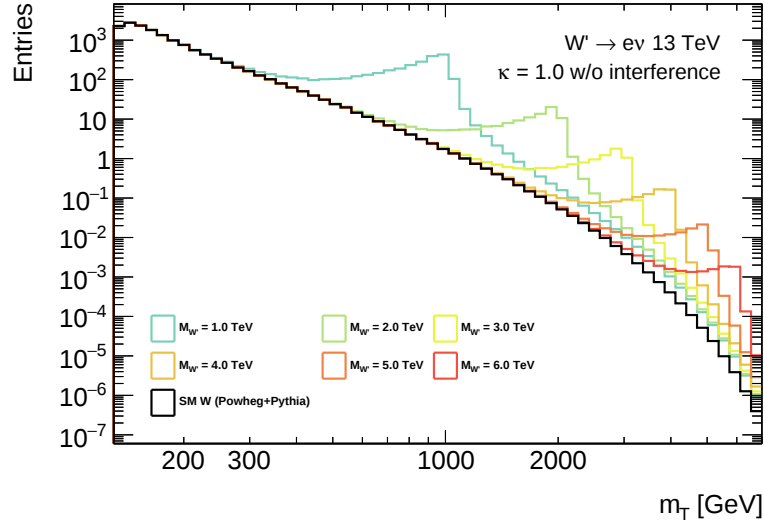


Figure 6.5.: Signal template in the electron channel without interference at $\kappa = 1.0$ for different pole masses of the W' . The distributions are scaled to an arbitrary integrated luminosity of 1.0 fb^{-1} .

When producing the template, this means that the content of bins can actually be negative, as can be seen in Equation (6.8), where the mixing-term might be negative.

$$|\mathcal{M}|^2 - |\mathcal{M}_{\text{SM}}|^2 = |\mathcal{M}_{\text{SSM}}|^2 + 2 \text{Re}(\mathcal{M}_{\text{SM}}^* \cdot \mathcal{M}_{\text{SSM}}) \quad (6.8)$$

Using the same technique as described in Section 6.1.1, an invariant mass dependent weight [55] can be derived which includes this mixing-term, and factorizes from the weight determined before. The additional factor which needs to be included to introduce the interference effect via the mixing-term is

$$w_{\text{int}} = 1 + 2 \frac{(\hat{s} - M_W^2)(\hat{s} - M_{W'}^2)/\kappa + \hat{s}^2(\Gamma_W/M_W)(\Gamma_{W'}/M_{W'})}{(\hat{s} - M_W^2)^2 + (\hat{s}\Gamma_W/M_W)^2} \quad (6.9)$$

where $\hat{s} = m_{\ell\nu}^2$, and $\Gamma_{W'}$ is determined as stated in Section 6.1.1. The total weight is then $w \times w_{\text{int}}$. Using this weight, the corresponding signal templates can be produced, as seen in Figure 6.6. There it becomes apparent, that at transverse masses below the pole mass, the expected number of events in the signal template is indeed negative. This results in a decrease in the event yield in the distribution containing both signal and background. The magnitude of this decrease scales with κ , as is expected since the weight depends on κ .

In Figure 6.7 the contribution from the W background has been added to the template, to visualize the destructive interference. The templates including interference retain the scaling of the resonance width and the cross-section with κ from their non-interference counterparts. Figure 6.7b shows negative bin contents for bins below the resonance peak, which Figure 6.7a does not. As the expected number of events from the off-shell W drops, the relative impact of the destructive interference effect becomes larger closer to the W' resonance region. Nevertheless, a negative number of events in the distribution of the signal plus the background is unphysical. This effect is most probably an artifact of the different corrections being applied to the W and the W' sample (NNLO QCD + NLO EW for

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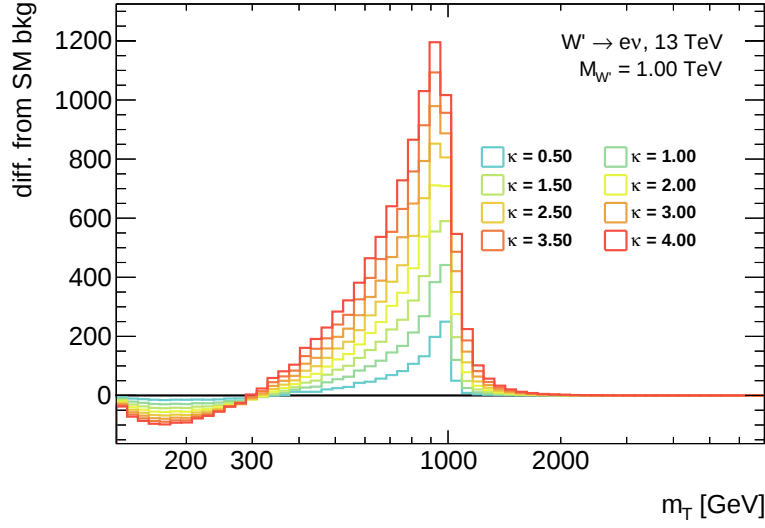


Figure 6.6.: Transverse mass distributions in the electron channel for a W' pole mass of $M_{W'} = 1$ TeV, excluding the SM W contribution (POWHEG+PYTHIA8) but including interference, as a function of different values of κ . The distributions are scaled to an arbitrary integrated luminosity of 1 fb^{-1} .

the background as opposed to NNLO QCD + LO EW for the signal), which was already mentioned in the previous section. This can be studied by calculating an additional weight variant, which includes interference effects, but also directly includes the SM W background. This corresponds to

$$w_{\text{int},2} \propto |\mathcal{M}|^2 = |\mathcal{M}_{\text{SM}}|^2 + |\mathcal{M}_{\text{SSM}}|^2 + 2 \text{Re}(\mathcal{M}_{\text{SM}}^* \cdot \mathcal{M}_{\text{SSM}}) \quad (6.10)$$

$$w_{\text{int},2} = \frac{(2\hat{s} - M_W^2 - M_{W'}^2)^2 + \hat{s}^2(\Gamma_W/M_W + \Gamma_{W'}/M_{W'})^2}{(\hat{s} - M_W^2)^2 + (\hat{s}\Gamma_W/M_W)^2} \quad (6.11)$$

Using this weight instead of w_{int} (Equation (6.9)) results in an equivalent shape as in Figure 6.7. This can be seen in Figure E.1, where no negative bin contents are visible. In Figure 6.8 the signal template can be seen for fixed $\kappa = 1.0$ and different pole masses. Also here, the impact of the destructive interference increases towards the high end of the spectrum, as the W off-shell contribution decreases.

Corresponding distributions for the signal templates determined in the muon channel can be found in Appendix A.2. As is the case without interference, Figure A.4 shows the lower muon momentum resolution.

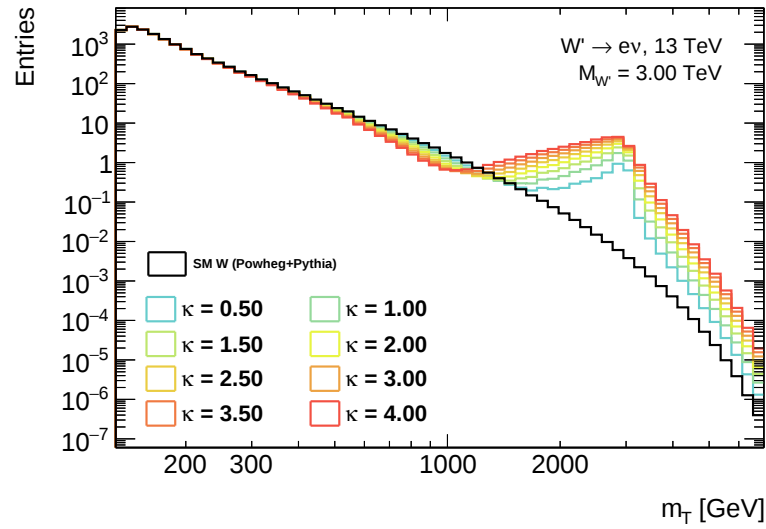
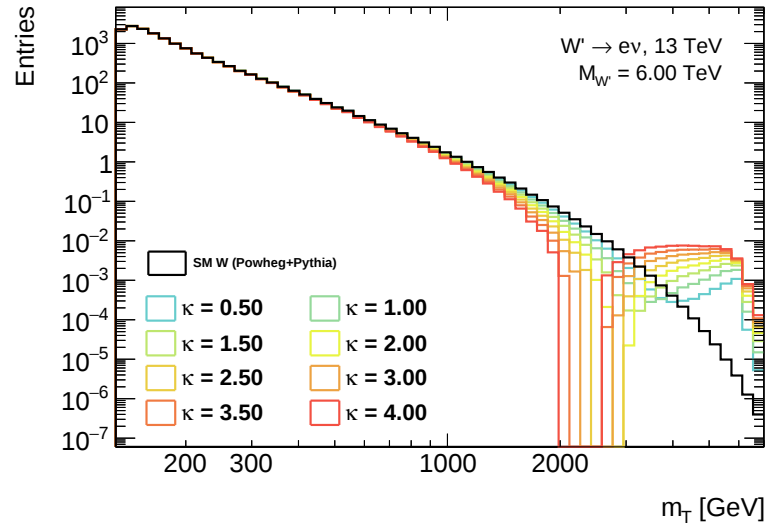
(a) $M_{W'} = 3 \text{ TeV}$ (b) $M_{W'} = 6 \text{ TeV}$

Figure 6.7.: Transverse mass distributions in the electron channel for four different W' pole masses, added to the SM W contribution (POWHEG+PYTHIA8) and including interference. The distributions are shown for different values of κ and are scaled to an arbitrary integrated luminosity of 1 fb^{-1} .

6.2. Bayesian statistics

The goal of the statistical analysis of data is to quantify the compatibility between the observed data, and the signal hypothesis being tested. There are two popular approaches to this problem, the frequentist and the Bayesian approach. This analysis makes use of the Bayesian approach, which in

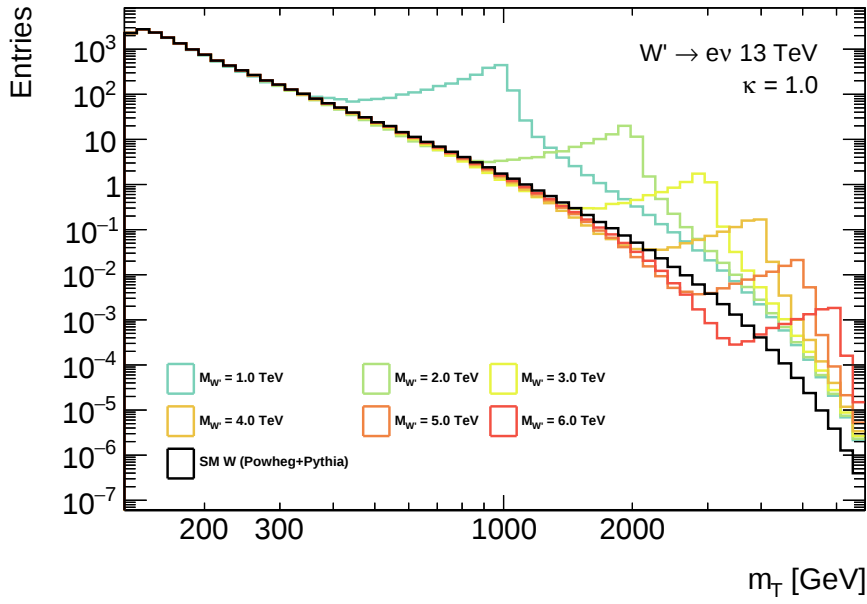


Figure 6.8.: Signal template in the electron channel with interference at $\kappa = 1.0$ for different pole masses of the W' . The distributions are scaled to an arbitrary integrated luminosity of 1 fb^{-1} .

turn is based on Bayes' theorem [56]:

$$P(A|B) = \frac{P(B|A)P(A)}{P(B)} \quad (6.12)$$

This theorem allows for an interpretation of the conditional probability stated above as a *degree of believe* of A , given B . The expression can be linked with the interpretation of the experiment, by associating A with the hypothesis under test. In the case of this analysis, A can be identified with the SSM W' hypothesis. B on the other hand, is linked to the data set. $P(A|B)$ is then the probability, that the hypothesis of new physics is true, given that the data set under study has been measured.

To perform the statistical analysis, a parameter of interest must be chosen, on which an upper limit can be set. The probability density for this parameter can then be determined, making use of the background and signal expectation, as well as the data. In published searches for W' particles, the most common parameter of interest is σB , which is the product of the cross-section of the signal process and the branching ratio for the decay into the observed final states. This is achieved by scaling the integral of the signal templates, and relating the change in the number of events to the cross-section. By comparing the upper limit to the cross-section predicted by theory, a mass exclusion limit can be calculated.

In contrast, to be able to compare mass exclusion limits obtained for scenarios including and excluding destructive interference, it is useful to set an upper limit on the coupling ratio parameter κ instead. κ affects the signal shape in both cases. The interpretation of $\kappa = 1$ corresponding to the SSM W' hypothesis, allows the calculation of a mass exclusion limit. The parameter scales the normalization of the signal, the decay width and the strength of the interference effect if it is included, as opposed to

only affecting the total normalization of the template as for σB . Equation (6.12) can be re-formulated as

$$p(\kappa, \vec{\theta} | \vec{n}) = \frac{P(\vec{n} | \kappa, \vec{\theta}) p(\kappa)}{P(\vec{n})} \quad (6.13)$$

where $\vec{n}$ is the set of data and $\vec{\theta}$ is the ensemble of *nuisance* parameters. These parameters describe, generally speaking, parameters which are not the parameter of interest. They can be used to incorporate systematic uncertainties into the statistical description. By introducing a nuisance parameter for each systematic uncertainty, the impact of it on the distributions can be modified by varying the value of the parameter. For this to work, systematic uncertainties must be given in the form of a relative shift of the background. $p(\kappa, \vec{\theta} | \vec{n})$ is the probability density for a given value of κ and a specific set of nuisance parameters $\vec{\theta}$.

To arrive at an expression for this probability density, $P(\vec{n} | \kappa, \vec{\theta})$ needs to be defined. This can be done by making use of the *likelihood* $\mathcal{L}$ between the histograms for the background, the signal template and the data. By including the signal template, the likelihood introduces the required κ dependence in the probability density above. The likelihood can be written as

$$\mathcal{L}(\kappa, \vec{\theta}) = \prod_{k=1}^{N_{\text{chan}}} \prod_{l=1}^{N_{\text{bin}}} \text{Pois}(\lambda_{kl}(\kappa, \vec{\theta}), n_{kl}) \quad (6.14)$$

where n_{kl} is the number of data events in bin l and decay channel k and N_{chan} and N_{bin} are the number of channels and number of bins in the histograms, respectively. By changing whether the likelihood includes the electron channel histograms, the muon channel histograms, or both channels, a limit can be calculated for the channels in isolation, or a combined limit can be calculated. $\lambda_{kl}(\kappa, \vec{\theta}) = s_{kl}(\kappa) + b_{kl}(\vec{\theta})$ is the sum of the number of signal and background events in the same bin and channel and

$$\text{Pois}(\lambda, k) = \frac{\lambda^k e^{-\lambda}}{k!} \quad (6.15)$$

is the poisson probability for k events and the expectation value λ . Note that the number of signal events s_{kl} depends on κ , whereas b_{kl} does not. $\lambda_{kl}(\kappa, \vec{\theta})$ also has a dependence on the vector $\vec{\theta} = (\theta_1, \theta_2, \dots, \theta_i)$. $P(\vec{n})$ is the prior probability of the data set, which is taken as a normalization constant in the following, and does not have a specific meaning.

$p(\kappa)$ is the prior probability density for different values of κ and describes the knowledge about the parameter of interest, before taking the data. The final result is sensitive to changes of this prior probability density. Typically, an uninformative prior approach is chosen, which restricts the input to as little information as possible. Thus, the obvious choice would be a flat probability density, meaning that all values of κ are equally likely. This, however, leads to a problem when trying to make a comparison to results from calculating upper limits on σB , where interference effects are neglected. In that case, the variation of the signal template is done through the normalization A , where higher values of σB increase the number of events expected from the signal process, thus $A \propto \sigma B$. If the impact of κ on the width of the W' resonance, as well as the destructive interference is neglected, the normalization A of the signal template would be proportional to κ^2 , as can be seen from Equation (6.3), resulting in $A \propto \kappa^2$.

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As a consequence it follows that $\kappa^2 \propto \sigma B$. In the σB limit approach, a flat prior probability density is chosen for σB . Thus, in a direct comparison, the prior probability density has to be flat in κ^2 , rather than in κ in order to be consistent. A probability density $p(\kappa) = \kappa$ results in such a flat distribution in κ^2 . This however, represents a fundamentally different choice than the non-informative flat prior probability density in κ . Results featuring both prior choices will be shown in the following.

To isolate the probability density as a function of κ only, the dependence on nuisance parameters needs to be removed. This step is called marginalization. During the integration, the prior probability densities of the nuisance parameters $\phi(\theta_i)$ are included, which are chosen to be normally distributed. Using the likelihood $\mathcal{L}(\kappa, \vec{\theta})$ the probability density can be written as

$$p(\kappa, \vec{\theta} | \vec{n}) = \frac{\mathcal{L}(\kappa, \vec{\theta}) p(\kappa)}{P(\vec{n})} \quad (6.16)$$

and then integrated over $\vec{\theta}$ to arrive at

$$p(\kappa | \vec{n}) = \int p(\kappa, \vec{\theta} | \vec{n}) d\vec{\theta} = N \int \prod_{k=1}^{N_{\text{chan}}} \prod_{l=1}^{N_{\text{bin}}} \text{Pois}(\lambda_{kl}(\kappa, \vec{\theta}), n_{kl}) \prod_{i=1}^{N_{\text{sys}}} \Phi(\theta_i) d\vec{\theta} \quad (6.17)$$

where N can be determined by requiring

$$\int_0^\infty p(\kappa | \vec{n}) d\kappa = 1. \quad (6.18)$$

Subsequently, integration over the obtained probability density yields the 95 % confidence level (CL) limit of κ_{up} .

$$\int_{\kappa_{\text{up}}}^\infty p(\kappa | \vec{n}) d\kappa = 1 - 0.95 \quad (6.19)$$

A good way to assess if the result of the limit setting procedure is stable and reasonable, without looking at actual data and running a risk of optimizing with a bias, the algorithm can be run with the background sample as a stand in for the data. This expected limit corresponds to the result that should be obtained if the data is compatible with the background only hypothesis. The result of the calculation using the exact value of the background sample is called the *Asimov* limit [57]. Additionally, by varying the content of each bin of the background histogram according to a poisson distribution with the expectation value set to the bin content, the sensitivity of the upper limit to statistical fluctuations of the samples can be studied. Each random variation of the nominal bin content is called a *pseudo-experiment*, as it corresponds to repeating the measurement and obtaining a compatible result within the statistical uncertainty. This expected limit will be presented in the following sections along with the Asimov limit.

The implementation of the statistical procedure is using the Bayesian Analysis Toolkit (BAT) [58]. The treatment is based on software written for Z' coupling limits [59], but was heavily modified for this analysis.

BAT runs a minimization of the negative logarithm of $\mathcal{L}$ using the Metropolis algorithm [60] to find the global minimum. Subsequently, using MINUIT [61], the minimal mode is found. BAT needs to be able to vary the parameter of interest as well as the nuisance parameters continuously for this likelihood

optimization. Since the nuisance parameters simply scale the provided systematic shifts, this is not a problem. The signal templates, on the other hand, have been reweighted to discrete values of κ , meaning that intermediate steps are not readily available. To overcome this, the κ distribution for each bin in m_T is interpolated. The resulting set of functions is used like a histogram, where the effective m_T bin content can be determined from the value of the function at that point, providing BAT with satisfactory input. After completion of the optimization step, the 95 % CL value κ_{up} can be

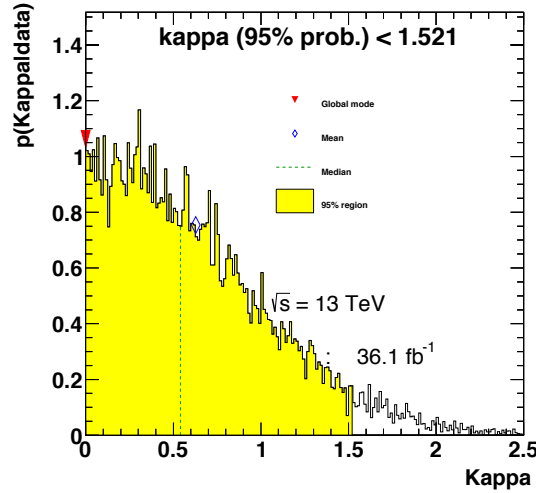


Figure 6.9.: Parameter probability distribution for the parameter of interest κ as determined by BAT, for the Asimov limit in the electron channel, at a pole mass of $M_{W'} = 6$ TeV.

extracted directly from the distribution of κ values generated during the algorithmic procedure. Such a distribution is shown in Figure 6.9 as an example for a pole mass of $M_{W'} = 6$ TeV for the Asimov limit in the electron channel. The area of the distribution up to the upper limit κ_{up} is highlighted. For the expected limit, a variation of the background expectation within statistical uncertainties is performed. For each pseudo-experiment, this results in an equivalent distribution as the one seen in Figure 6.9. The ensemble of κ_{up} is then used to determine the median, as well as the standard deviation, and the 95 % quantile. A distribution of κ_{up} values in the ensemble is shown in Figure 6.10 for the electron channel and 1000 pseudo-experiments and a pole mass of 6 TeV. The $\pm 1\sigma$ and $\pm 2\sigma$ bands are also shown.

6.3. Non-interference limits

By using 118 templates which are produced for various pole masses only including the weight described in Section 6.1.1, upper limits on κ neglecting interference effects between the W' and the W are calculated for the electron or muon channel, or by using a combination of both. The pole masses

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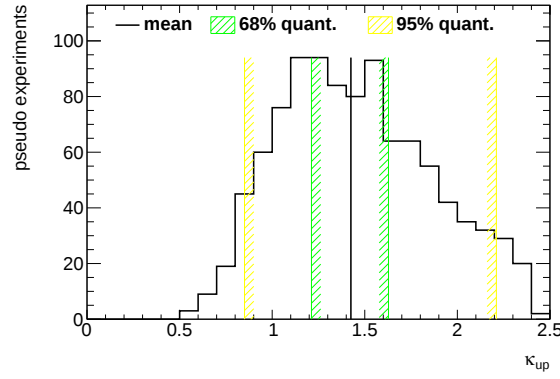


Figure 6.10.: Distribution of the 95 % CL upper limits for the ensemble of pseudo-experiments for the electron channel at a pole mass of $M_{W'} = 6$ TeV.

range from 150 GeV to 6 TeV in steps of 50 GeV. The parameter range for κ is $[10^{-6}, 2.5]$. Values for κ are equidistant on a logarithmic scale, in order to improve the interpolation at low values. The expected limits are calculated using 500 pseudo-experiments.

The upper limit κ_{up} is shown as a function of the pole mass $M_{W'}$ in Figures 6.11 to 6.13 for electrons, muons and both electrons and muons combined. Additionally, limits are calculated and shown for the constant and linear prior probability density choice for κ . In all cases, the upper limit curves are compatible with no coupling for low values of m_T , only at around $M_{W'} = 3$ TeV do the upper limits deviate from zero.

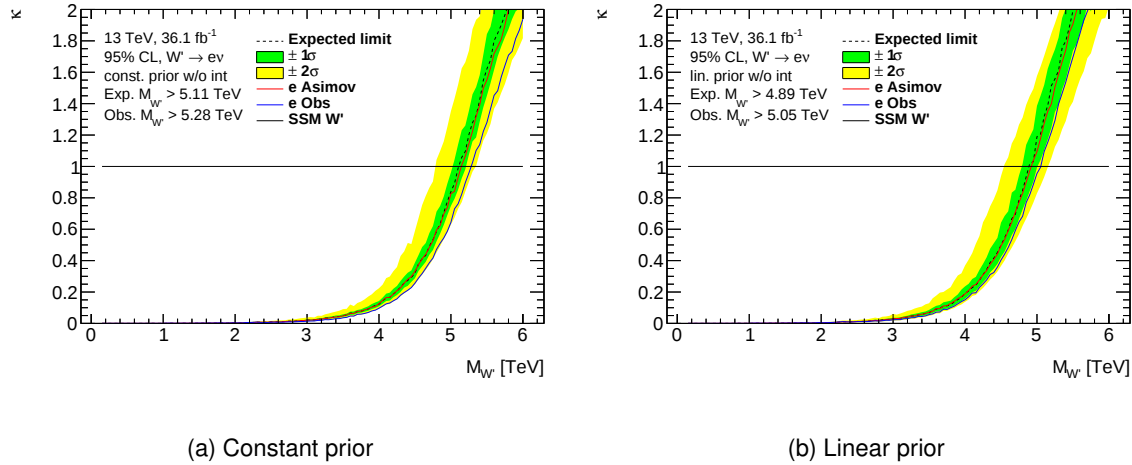


Figure 6.11.: 95 % CL upper limits on κ in the electron channel with a constant (a) and linear (b) prior and without interference.

In the electron channel shown in Figure 6.11, the Asimov limit is very close to the median value of the expected limit. Aside from the median value, the values for $\pm 1\sigma$ and $\pm 2\sigma$ are drawn as bands in different colors, respectively. These bands give an indication of the stability of the expected limit with respect to statistical fluctuations of the pseudo-data. The observed limit from data is consistently stronger than the expected limit. A horizontal line indicates $\kappa = 1$. This corresponds to the SSM

which postulates a W' boson with equal coupling to fermions as the W . As a consequence, the SSM W' can be excluded for masses

$$M_{W'} < M_{W'}|_{\kappa_{\text{up}}=1} \equiv M_{W',\text{min}}, \quad (6.20)$$

since the (pseudo-)data constrains the value of κ to be smaller than unity. This is then incompatible with the hypothesis. Consequently, a mass exclusion limit $M_{W',\text{min}}$ can be extracted by finding the intersection between the limit curve and unity. For the electron channel alone, expected mass exclusion values of 5.11 TeV (constant prior, Figure 6.11a) and 4.89 TeV (linear prior, Figure 6.11b) can be found at 95 % CL. The corresponding observed limits are 5.28 TeV (constant prior, Figure 6.11a) and 5.05 TeV (linear prior, Figure 6.11b).

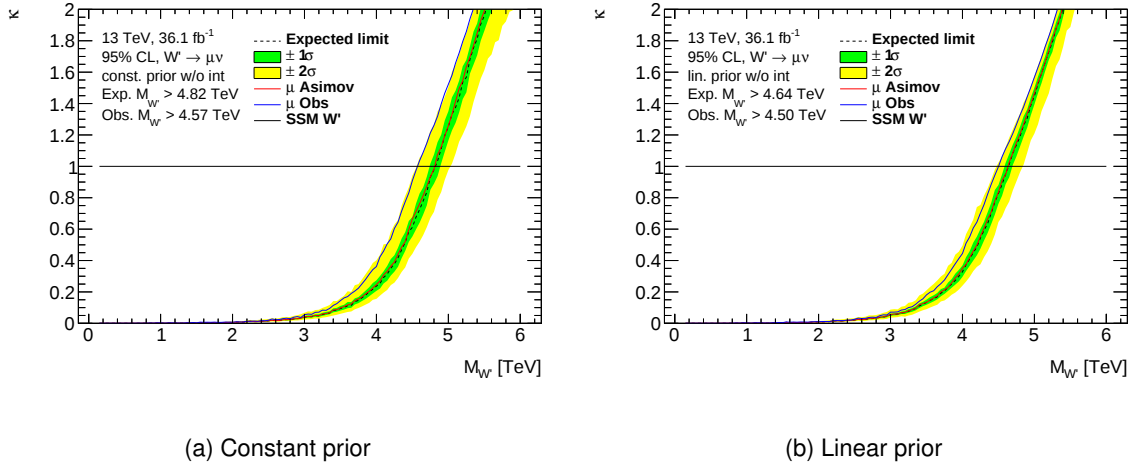


Figure 6.12.: 95 % CL upper limits on κ in the muon channel with a constant (a) and linear (b) κ prior and without interference.

For the muon channel in Figure 6.12 the expected limit curve departs from zero slightly under $M_{W'} = 3$ TeV, and then rises. The Asimov limit also appears compatible with the median of the expected limit. By finding the intersection with $\kappa = 1$, expected mass exclusion limits of $M_{W'} > 4.82$ TeV (constant prior, Figure 6.12a) and $M_{W'} > 4.64$ TeV (linear prior, Figure 6.12b) can be found at 95 % CL for the prior choices. In this case, the observed limit is consistently weaker than the expected limit. The observed mass exclusion limits are $M_{W'} > 4.57$ TeV (constant prior, Figure 6.12a) and $M_{W'} > 4.50$ TeV (linear prior, Figure 6.12b).

Figure 6.13 shows the combined limit obtained using both the electron and the muon channel. The general appearance of the limit curve is very similar to the ones seen for electrons and muons in isolation. By using both channels simultaneously to assess compatibility, stronger expected mass exclusion limits of $M_{W'} > 5.27$ TeV (constant prior, Figure 6.13a) and $M_{W'} > 5.06$ TeV (linear prior, Figure 6.13b) can be found at 95 % CL for both prior choices. The corresponding observed mass exclusion limits are $M_{W'} > 5.15$ TeV (constant prior, Figure 6.13a) and $M_{W'} > 5.00$ TeV (linear prior, Figure 6.13b).

6. Statistical analysis

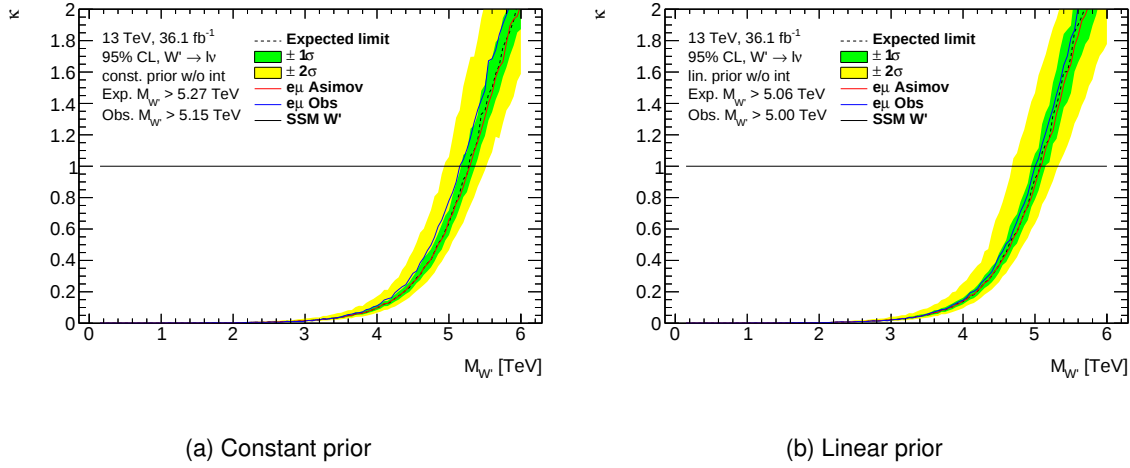


Figure 6.13.: 95 % CL upper limits on κ in the combination of the electron and the muon channel with a constant (a) and linear (b) κ prior and without interference.

6.4. Interference limits

Using the signal templates which are produced with the additional interference weight, the same limit setting procedure can be conducted. Technically, no different behavior needs to be implemented, as the fitted functions which reproduce the m_T bin contents as continuous functions of κ include all the information on the destructive interference effect.

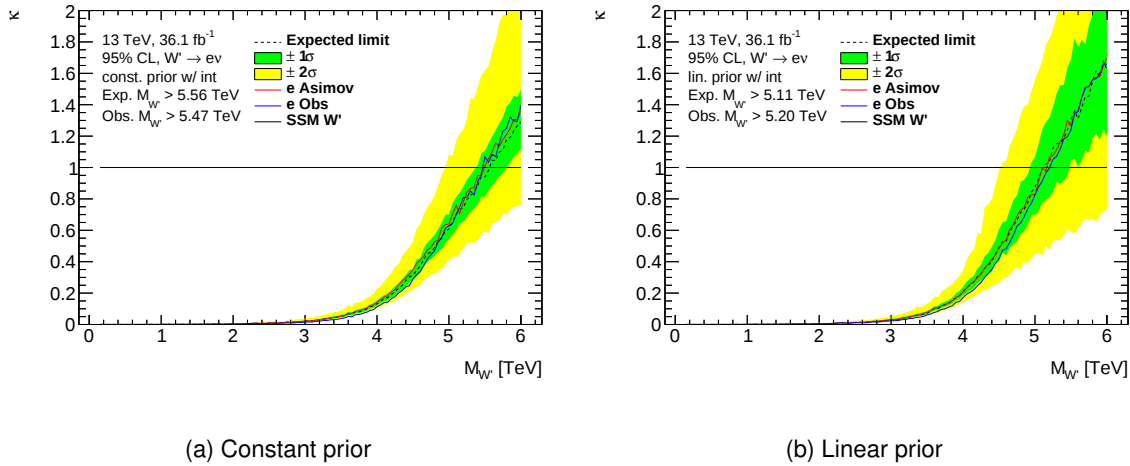


Figure 6.14.: 95 % CL upper limits on κ in the electron channel with a constant (a) and linear (b) κ prior and including interference.

Analogously to what was seen before, Figure 6.14 shows the expected, Asimov and observed limit curves in the electron channel as a function of $M_{W'}$ for the constant and linear prior choice, incorporating destructive interference. Again, the intersection between the $\kappa = 1$ yields the mass inclusion limit, the interpretation of unity representing the SSM W' hypothesis remains unchanged. These intersections are found to be at 5.56 TeV (constant prior, Figure 6.14a) and 5.11 TeV (linear prior, Fig-

ure 6.14b) at 95 % CL for the expected limit. The Asimov limit is compatible with the median value of the expected limit, the observed limit is marginally stronger. The corresponding mass exclusion limits are 5.47 TeV (constant prior, Figure 6.14a) and 5.20 TeV (linear prior, Figure 6.14b).

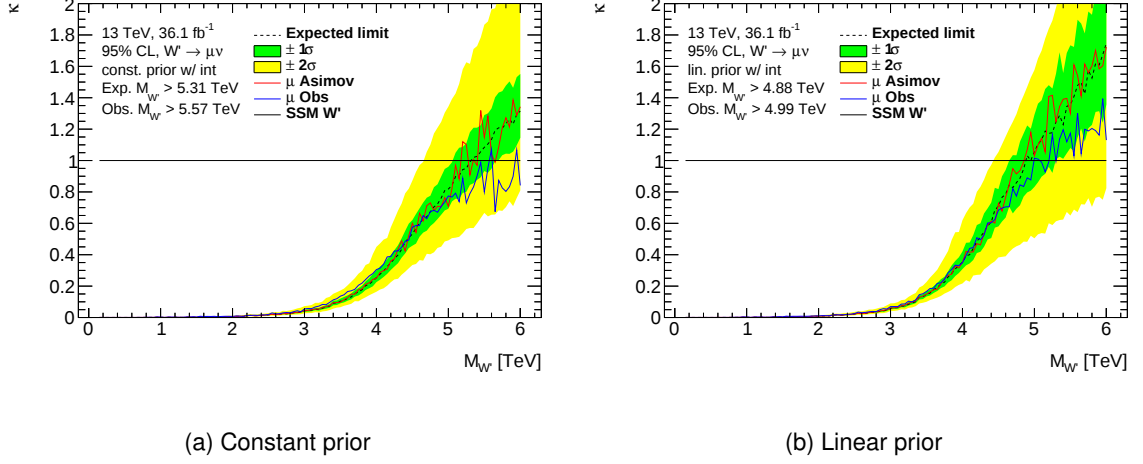


Figure 6.15.: 95 % CL upper limits on κ in the muon channel with a constant (a) and linear (b) κ prior and including interference.

Figure 6.15 shows the equivalent limit curves for the muon channel. Here, mass exclusion limits of $M_{W'} > 5.31$ TeV (constant prior, Figure 6.15a) and $M_{W'} > 4.88$ TeV (linear prior, Figure 6.15b) are found for both prior scenarios at 95 % CL for the expected limit. The corresponding observed mass exclusion limits are $M_{W'} > 5.57$ TeV (constant prior, Figure 6.15a) and $M_{W'} > 4.99$ TeV (linear prior, Figure 6.15b). It is notable, that the observed limit curve appears to depart from the expected and Asimov limit curves at around 5 TeV for both priors. This might be related to the shape of the muon signal templates for the interference case. Here, the data appears to be more compatible with the signal expectation than the background alone. At high masses, the signal resonance for the muon channel becomes less pronounced and also begins to fall outside of the histogram used for the statistical treatment (see Figure A.6).

The combined limit curves seen in Figure 6.16 are stronger than either electron or muon channel on its own. The limits are consistent with what was seen in the non-interference results. Expected mass exclusion limits of $M_{W'} > 5.96$ TeV (constant prior, Figure 6.16a) and $M_{W'} > 5.51$ TeV (linear prior, Figure 6.16b) are found for the two prior choices at 95 % CL. For the constant prior case, the observed upper limit κ_{up} does not cross unity within the pole mass range up to 6 TeV. Hence, no mass exclusion limit can be extracted for this case. This effect is related to the deviation of the observed limit curve from the expected one that is seen in the muon channel. In the linear prior case, the observed mass exclusion limit is $M_{W'} > 5.78$ TeV.

6. Statistical analysis

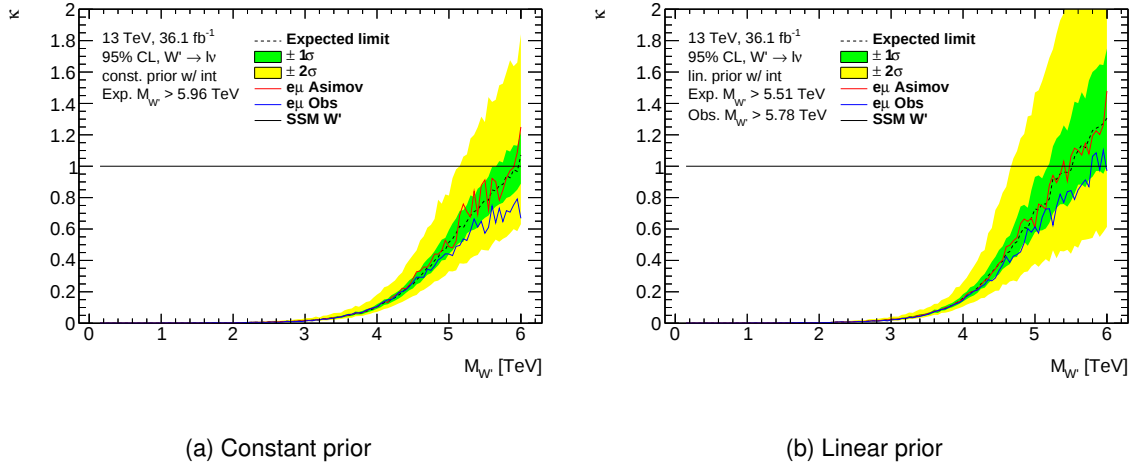


Figure 6.16.: 95 % CL upper limits on κ in the combination of the electron and the muon channel with a constant (a) and linear (b) κ prior and including interference.

6.5. Comparison of results

The analysis published in [50] uses a limit setting approach with σB as a parameter of interest and neglects interference effects. The compatibility of the limit setting approach using κ with the results obtained from the σB variant can be tested. For this, the κ limit setting procedure is slightly modified to approximate the σB approach. This is done by dropping the factor of κ that affects the resonance width (see Equation (6.6)) and only including the non-interference part of the weight (see Equation (6.3)). In this configuration, the only impact that the parameter of interest has on the signal template is the total normalization. As mentioned in Section 6.2, a linear prior is chosen so that the setup is consistent. The signal templates, limit curves and mass exclusion limits for this variant can be found in Appendix C. The σB limit setup is executed with the data, background and signal templates used in this analysis. The corresponding limit curves can be found in Appendix D. They are in general agreement with the results from the published analysis in [50]. The muon channel shows a moderate discrepancy, which is likely to originate from the simplified treatment for this analysis.

Table 6.1.: Observed and expected mass exclusion limits from upper limits on κ and σB at 95 % CL. Interference effects are neglected, and the decay width does not depend on κ .

upper limit on:	σB		$\kappa, \Gamma \neq \Gamma(\kappa)$	
channel	expected [TeV]	observed [TeV]	expected [TeV]	observed [TeV]
electron	5.10	5.22	5.10	5.22
muon	4.76	4.52	4.74	4.50
combined	5.22	5.12	5.23	5.12

Since the upper limits on κ and σB cannot be compared directly, a comparison of the mass exclusion limits can be made. The mass exclusion limits obtained using these two scenarios are shown in Table 6.1 and are compatible with each other. Minor discrepancies are most likely to originate from small technical differences in the limit setting algorithms. The upper limits on κ can thus be assumed

viable and compatible with the existing approach.

Table 6.2.: Comparison of expected mass exclusion limits from upper limits on κ in the various scenarios.

channel / type	expected mass exclusion limit at 95 % CL [TeV]				increase [GeV]	
	w/o interference		w/ interference			
	const. prior	lin. prior	const. prior	lin. prior	const. prior	lin. prior
electron	5.11	4.89	5.56	5.11	450	220
muon	4.82	4.64	5.31	4.88	490	240
combined	5.27	5.06	5.96	5.51	690	450

Table 6.2 shows the expected mass exclusion limits obtained from upper limits on κ for the constant and linear prior choices, as well as for included and excluded interference effects. Within each interference scenario and channel, the linear prior choice yields a weaker mass exclusion limit. This can be understood by considering the impact of the prior choice. A linear prior $p(\kappa) = \kappa$ enhances the likelihood for larger values of κ , thereby shifting the upper limit upwards.

The comparison of the two interference scenarios shows that including interference effects results in stronger expected mass exclusion limits. The signal shape differs significantly when interference is included. The limit setting becomes sensitive at masses below the resonance region, where the interference term still has an impact. Without interference, this region is affected by the signal to a negligible degree. Thus, the variation of the parameter of interest results in more significant changes in the background-plus-signal expectation. This should result in a stronger constraint on κ .

The mass exclusion limits calculated for the consistency check (Table 6.1) appear to be relatively close to the constant prior mass exclusion limits without interference effects (second column in Table 6.2). The linear prior scenario yields lower mass exclusion limits than the consistent limits, which also use a linear prior. This is most likely due to the variation of the decay width for the regular limits on κ . The constant prior is the most uninformative choice when including the impact of κ on the decay width. It is therefore reasonable to use this prior to determine the observed upper limits and mass exclusion limits.

Table 6.3.: Comparison of observed mass exclusion limits determined using a constant prior for κ .

channel	observed mass exclusion limit at 95 % CL [TeV]		increase [GeV]
	w/o interference	w/ interference	
electron	5.28	5.47	190
muon	4.57	5.57	1000
combined	5.15	$> 6.00^1$	≈ 850

Table 6.3 shows the observed mass exclusion limits using a constant prior, both with and without interference effects. As is the case for the expected mass exclusion limits, including interference effects leads to an increase in the observed mass exclusion limits. For the electron channel, the

¹The observed limit does not cross $\kappa = 1$ below a pole mass of 6 TeV, see Figure 6.16a.

6. Statistical analysis

expected mass exclusion limit increases by about 450 GeV, whereas the increase amounts to 190 GeV in the observed case. In the muon channel the expected mass exclusion limit increases by about 490 GeV. The observed mass exclusion limit in the muon channel increases quite drastically by about 1000 GeV. The large increase is connected to the deviation of the observed limit curve from the expected limit at around 5 TeV (see Figure 6.16a), that emerges in the muon channel when including interference effects. This is likely related to the shape of the signal templates for the muon channel, and specifically the impact of the interference weight on them. In the combination of both the electron and muon channel, no mass exclusion limit can be extracted, since the observed limit curve does not intersect with $\kappa = 1$ within the W' pole mass region up to 6 TeV. This corresponds to an observed mass exclusion limit of > 6 TeV, and would indicate an increase by about 850 GeV.

7. Summary and conclusion

A search for new heavy charged W' gauge bosons has been performed with data from the ATLAS detector amounting to 36.1 fb^{-1} of integrated luminosity at $\sqrt{s} = 13 \text{ TeV}$. The decay of the hypothesized W' to electrons and muons has been studied. The treatment of muons is lacking some features of the full published analysis in [50]. The decay results in a neutrino alongside the charged lepton. Using the reconstructed missing transverse momentum E_T^{miss} the kinematics of the neutrino are taken into account. Contributions from background processes are estimated using Monte Carlo simulations and a data driven approach is used for the Multijet background in the electron channel. The comparison between the background expectation and the observed data shows a reasonable compatibility and suggests a good modeling of the physics processes. Since no excess is seen in the transverse mass spectrum for the data, upper limits on various parameters can be calculated. To this end, several systematic uncertainties and their impact on the background estimates are studied. The published analysis sets upper limits on the product of cross section and branching ratio σB , whereas this analysis focuses on the relative coupling strength

$$\kappa = \left(\frac{g_{W'}}{g_W} \right)^2. \quad (7.1)$$

The limit setting is performed at 95 % CL using a Bayesian statistics approach. By using a reweighting technique, signal templates, which contain the expected deviation from the background under the signal hypothesis, can be created for various pole masses. An additional weight can be introduced to account for interference effects between the W' and the W . The impact on the signal is studied and limits using both approaches are calculated. By considering the Sequential Standard Model, which postulates $\kappa = 1$, a lower mass exclusion limit can be derived using the upper limits on κ . This result is susceptible to changes in the statistical analysis. Two choices for a prior probability for the parameter of interest are tested and compared. By excluding the impact of κ on the width of the signal resonance width, mass exclusion limits consistent with the ones obtained from σB can be calculated. Thus, upper limits on κ appear to be a viable strategy.

Using a constant prior for κ , expected (observed) mass exclusion limits of 5.11 TeV (5.28 TeV) are found for the electron channel, 4.82 TeV (4.57 TeV) are found for the muon channel and 5.27 TeV (5.15 TeV) are found for the combination without interference effects. When including interference effects, the mass exclusion limits increase to 5.56 TeV (5.47 TeV) for the electron channel, 5.31 TeV (5.57 TeV) for the muon channel and 5.96 TeV ($> 6 \text{ TeV}^1$) for the combination. The observed combined limit with interference effects does not cross the SSM expectation of $\kappa = 1$ below 6 TeV. An exact mass exclusion limit can thus not be extracted.

7. Summary and conclusion

The limit setting approach using the κ parameter yields stronger mass exclusion limits when including interference effects. In all studied scenarios and channels, the inclusion of interference effects in the signal expectation raises the mass exclusion limit by $\gtrsim 200$ GeV. The signal resonance influences a larger part of the transverse mass spectrum if interference is included as opposed to only affecting a window around the pole mass if it is excluded.

The results found in this analysis can further be improved upon. Using the full amount of data that will be available at the end of Run 2, the statistical uncertainties will significantly decrease. Other possibilities of improvement would be the use of Monte Carlo samples with higher statistics. Even though the samples for the DY background are binned in mass, the statistical uncertainties of the Monte Carlo exceed the ones found in data at lower masses. Additionally, the inclusion of mass binned samples for the Diboson and Top backgrounds could render the currently used extrapolation procedure unnecessary for them. As indicated by the mass exclusion limits presented above, an extension of the W' pole mass range used for the search could be useful as the limits approach 6 TeV.

A. Muon channel signal templates

A.1. Without interference

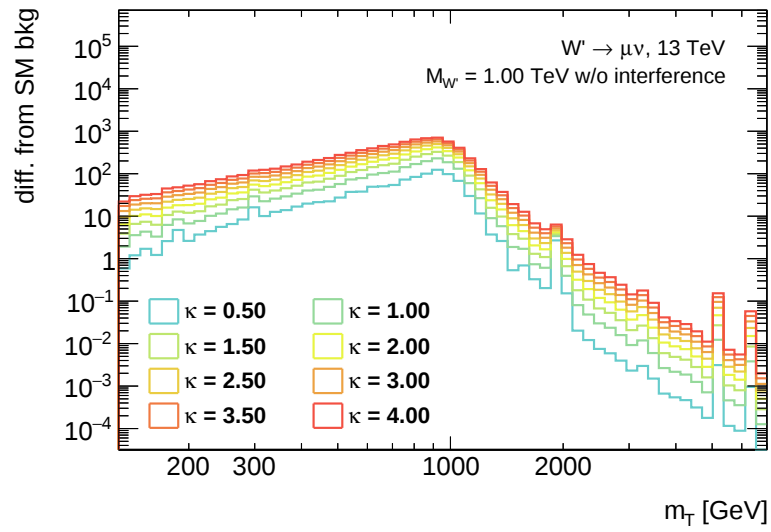
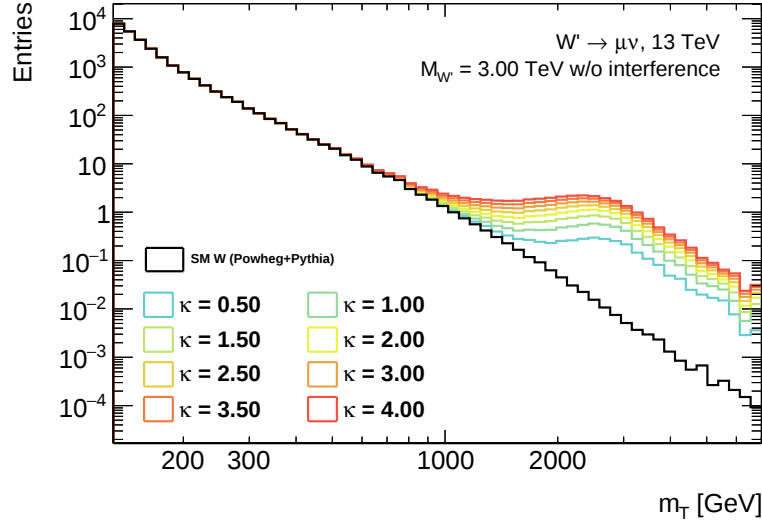
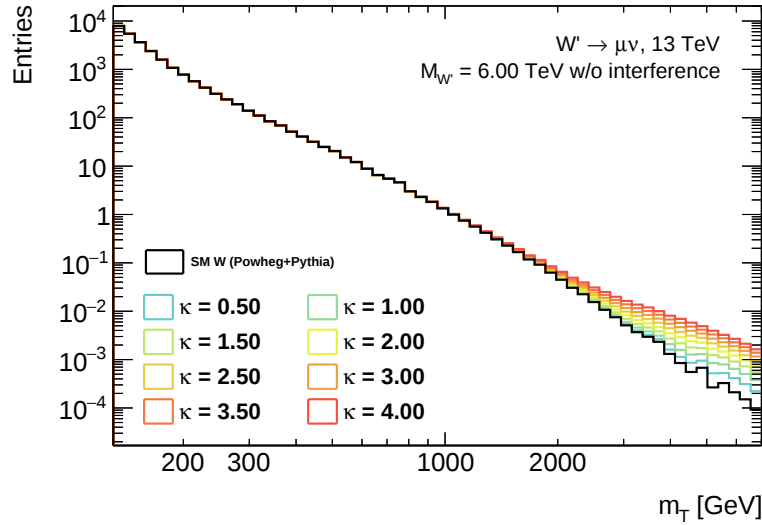


Figure A.1.: Transverse mass distribution in the muon channel for a W' pole mass of $M_{W'} = 1$ TeV, excluding the Standard Model contribution and interference effects, as a function of different values of κ . The distributions are scaled to an arbitrary integrated luminosity of 1 fb^{-1} .

A. Muon channel signal templates



(a) $M_{W'} = 3$ TeV



(b) $M_{W'} = 6$ TeV

Figure A.2.: Transverse mass distributions in the muon channel for different W' pole masses, added to the Standard Model W contribution (POWHEG+PYTHIA8) but excluding interference. The distributions show the results of different values of κ , indicated by the colors. The distributions are scaled to an arbitrary integrated luminosity of 1 fb^{-1} .

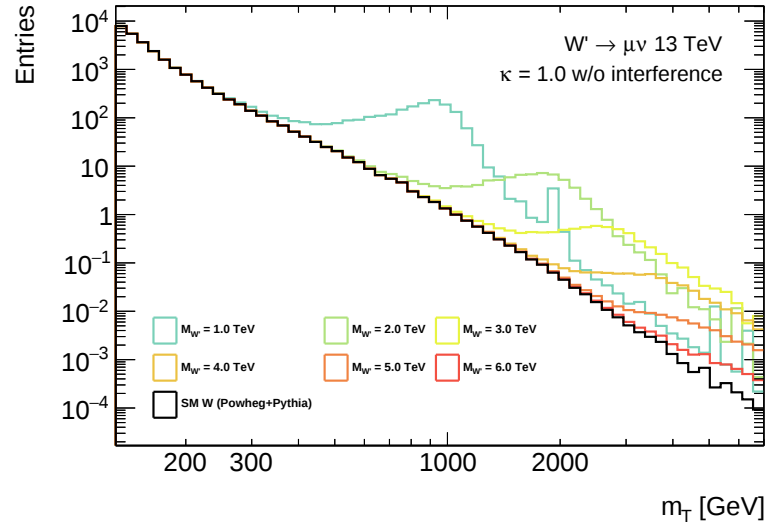


Figure A.3.: Signal template in the muon channel without interference at $\kappa = 1.0$ for different pole masses of the W' . The distributions are scaled to an arbitrary integrated luminosity of 1.0 fb^{-1} .

A.2. With interference

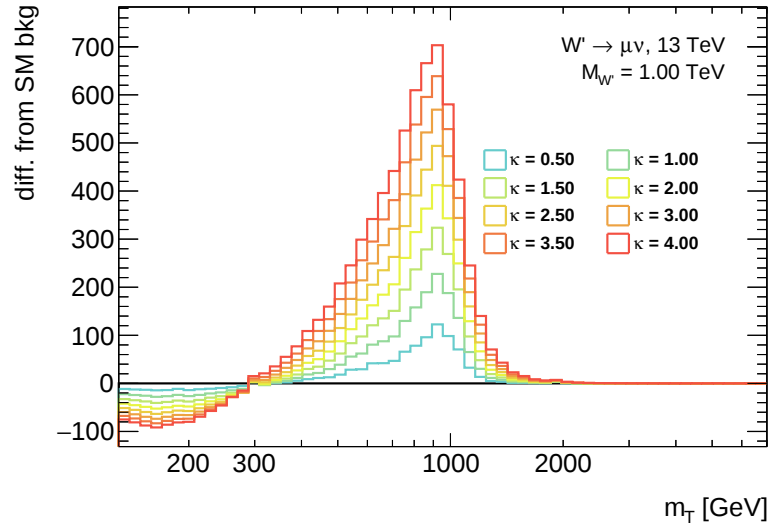
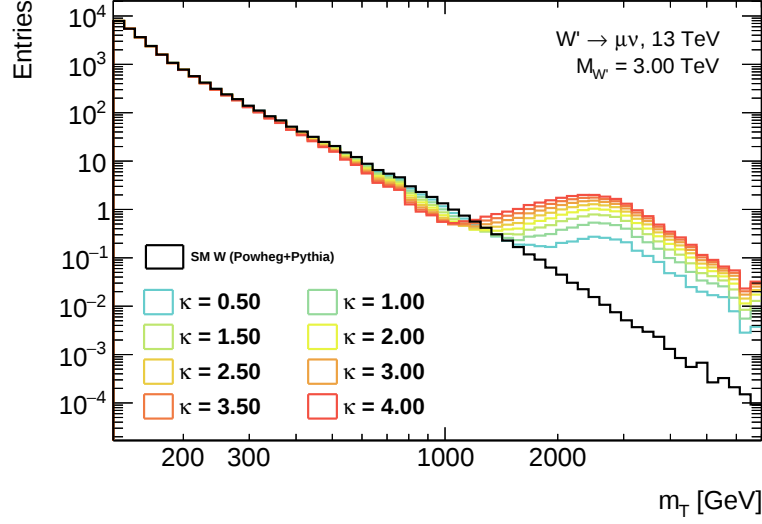
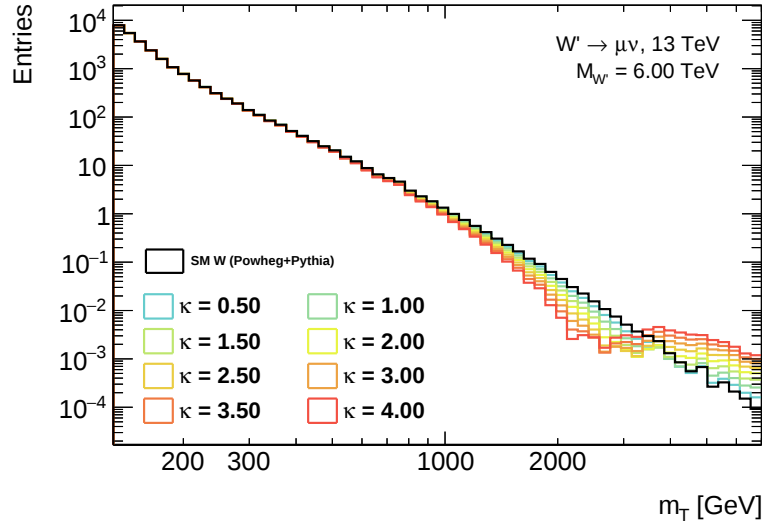


Figure A.4.: Transverse mass distributions in the muon channel for a W' pole mass of $M_{W'} = 1 \text{ TeV}$, excluding the Standard Model W contribution (POWHEG+PYTHIA8) but including interference, as a function of different values of κ . The distributions are scaled to an arbitrary integrated luminosity of 1 fb^{-1} .

A. Muon channel signal templates



(a) $M_{W'} = 3$ TeV



(b) $M_{W'} = 6$ TeV

Figure A.5.: Transverse mass distributions in the muon channel for four different W' pole masses, added to the Standard Model W contribution (POWHEG+PYTHIA8) and including interference. The distributions are shown for different values of κ and are scaled to an arbitrary integrated luminosity of 1 fb^{-1} .

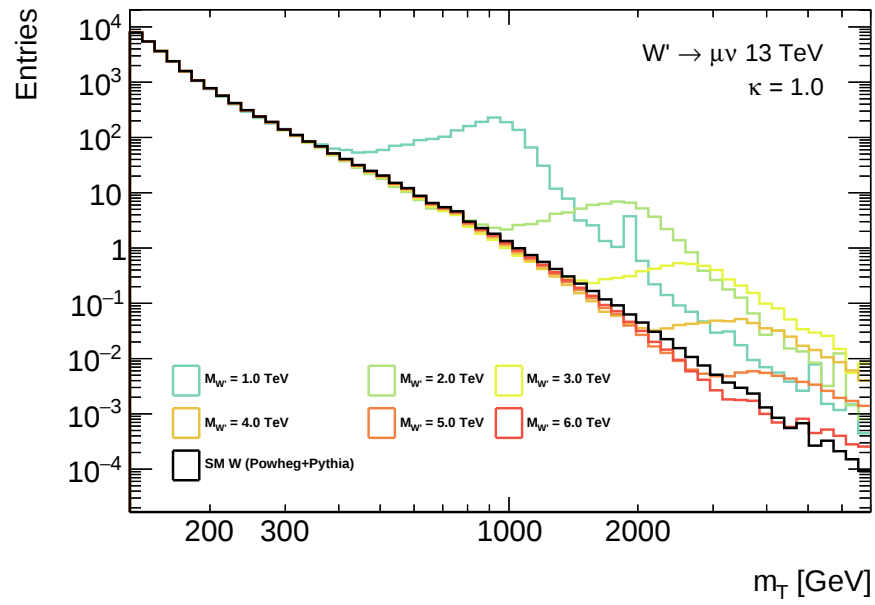
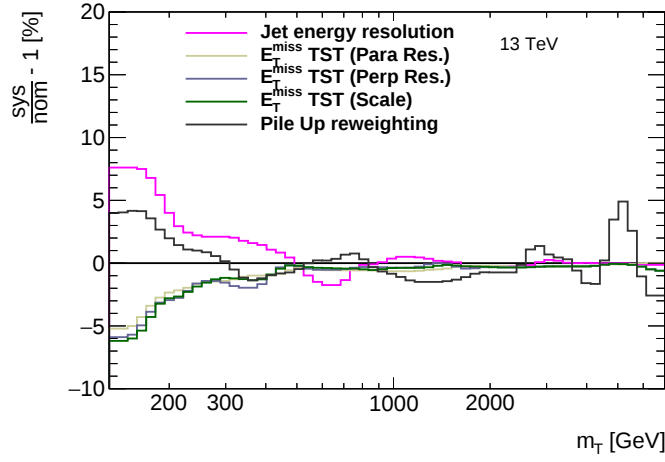
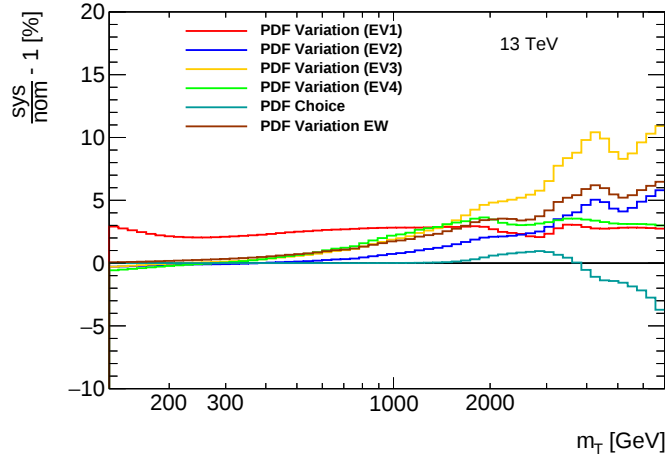


Figure A.6.: Signal template in the muon channel with interference at $\kappa = 1.0$ for different pole masses of the W'

A.3. Muon channel systematics



(a) Energy resolution and Pile-up uncertainties.



(b) Theoretical uncertainties.

Figure A.7.: Plots of the theory and energy resolution related systematic uncertainties which are considered for the analysis for the muon channel. The plots show the relative size of the deviation from the nominal background estimate in percent for different values of m_T .

B. Additional mass exclusion limits

Table B.1.: Comparison of expected and observed mass exclusion limits determined in the various scenarios.

channel / type	mass exclusion limit at 95 % CL [TeV]					
	w/o interference		w/ interference		increase [GeV]	
	const. prior	lin. prior	const. prior	lin. prior	const. prior	lin. prior
electron	5.11	4.89	5.56	5.11	450	220
electron observed	5.28	5.05	5.47	5.20	190	150
muon	4.82	4.64	5.31	4.88	490	240
muon observed	4.57	4.50	5.57	4.99	1000	490
combined	5.27	5.06	5.96	5.51	690	450
combined observed	5.15	5.00	—	5.78	—	780

C. Limits without variation of the resonance width

Upper limits on κ calculated using signal samples, where no interference weight is applied, and where the decay width Γ does not depend on Γ . The expected limits are calculated from 500 pseudo-experiments.

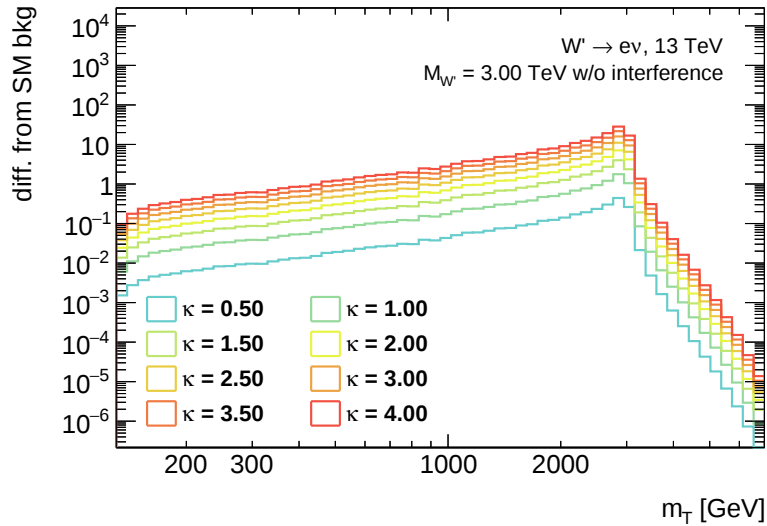


Figure C.1.: Signal template where Γ is independent of κ for the electron channel.

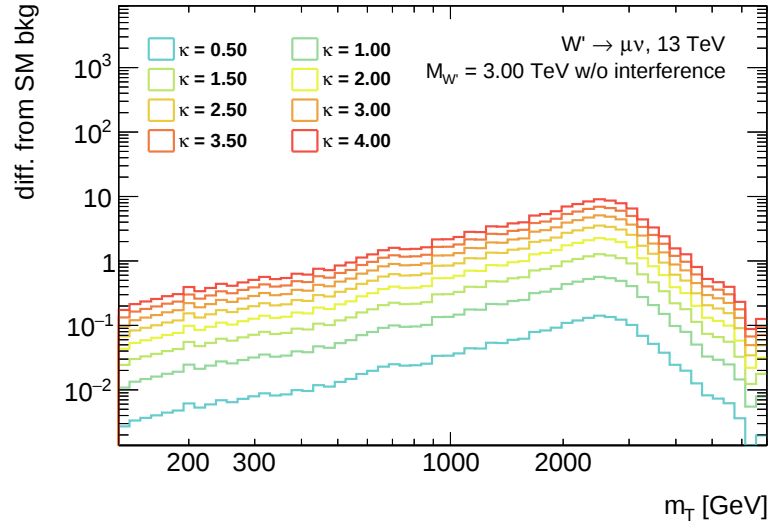


Figure C.2.: Signal template where Γ is independent of κ for the muon channel.

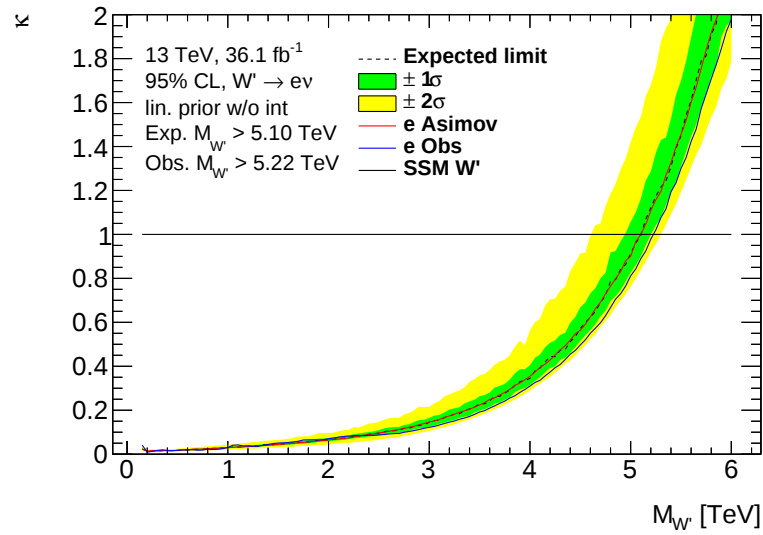


Figure C.3.: Upper limits on κ of the decay of $W' \rightarrow e\nu$.

C. Limits without variation of the resonance width

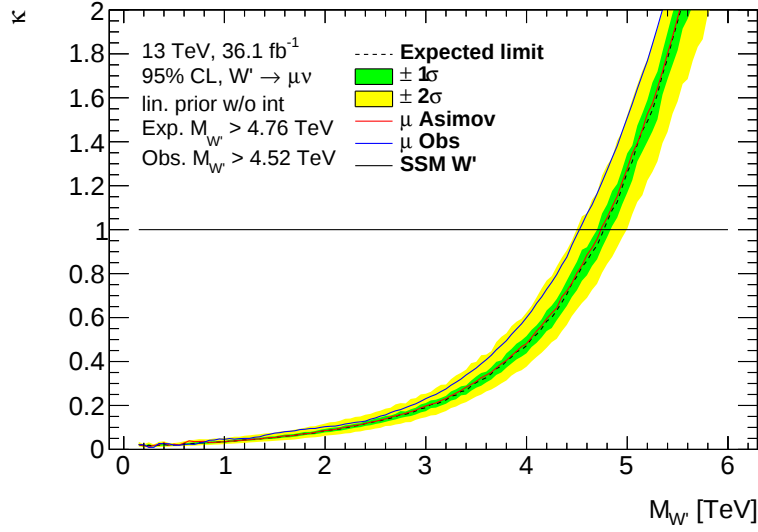


Figure C.4.: Upper limits on κ of the decay of $W' \rightarrow \mu\nu$.

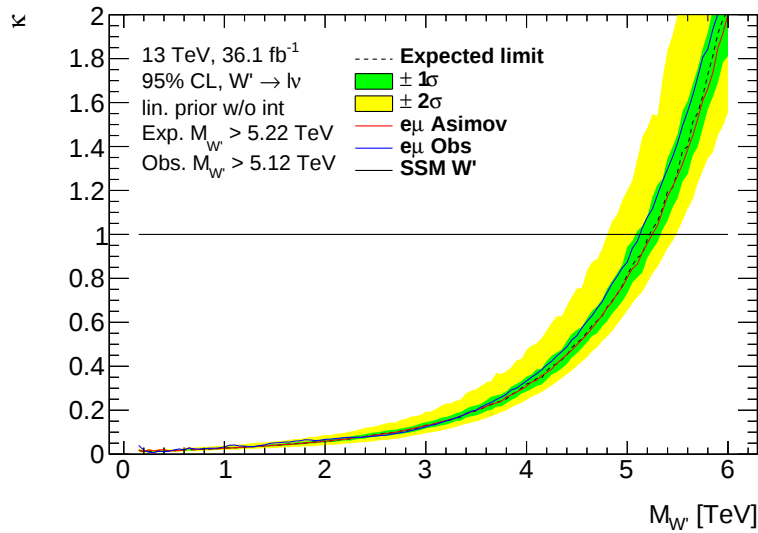


Figure C.5.: Upper limits on κ of the decay of $W' \rightarrow l\nu$.

D. Limits on cross section times branching ratio

Upper limits on σB of the W' are calculated as a comparison to reference [50]. The software is slightly modified from [62].

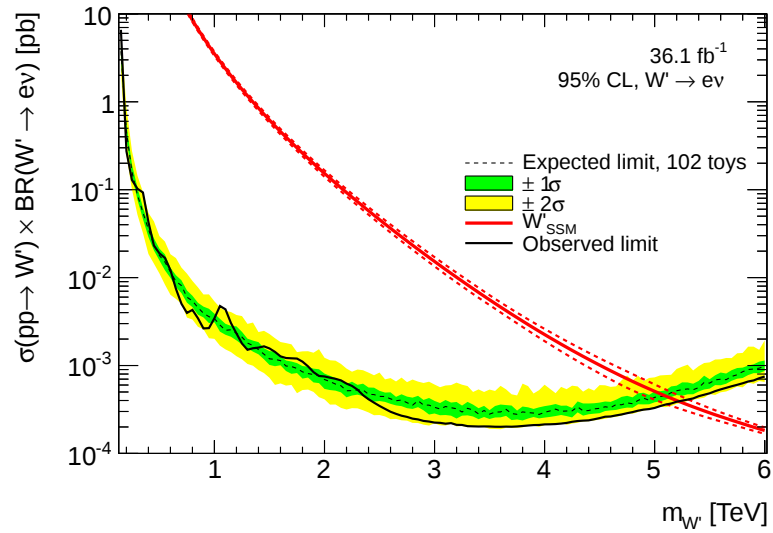


Figure D.1.: Upper limits on σB of the decay of $W' \rightarrow e\nu$.

D. Limits on cross section times branching ratio

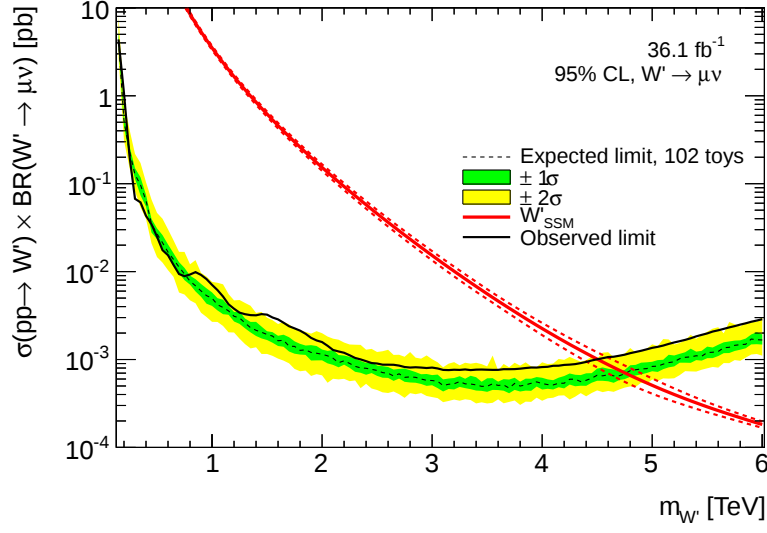


Figure D.2.: Upper limits on σB of the decay of $W' \rightarrow \mu\nu$.

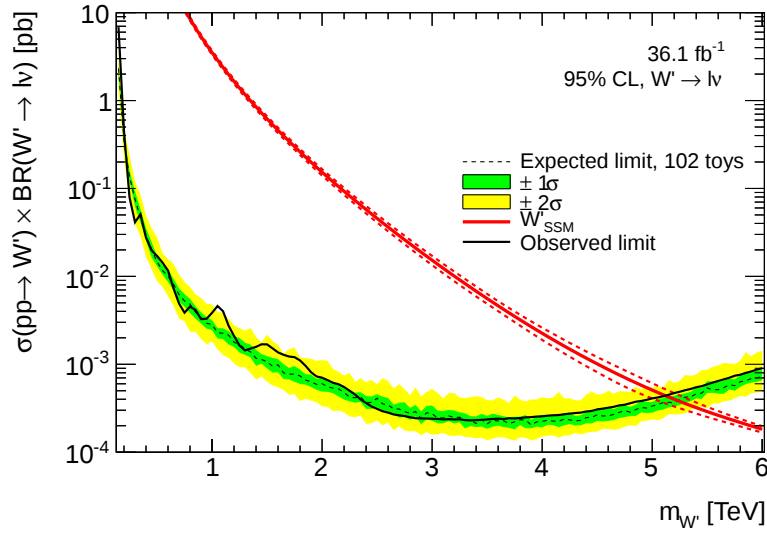
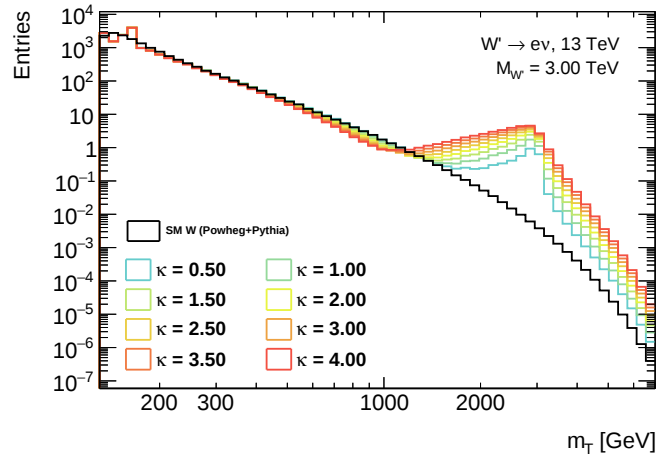
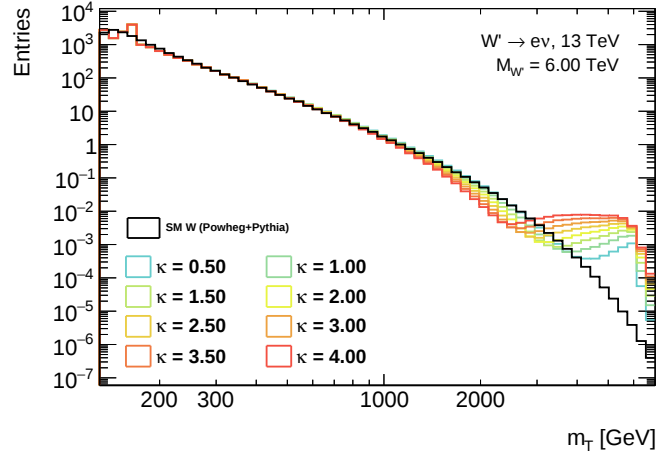


Figure D.3.: Upper limits on σB of the decay of $W' \rightarrow l\nu$.

E. Electron signal template plots including the SM background



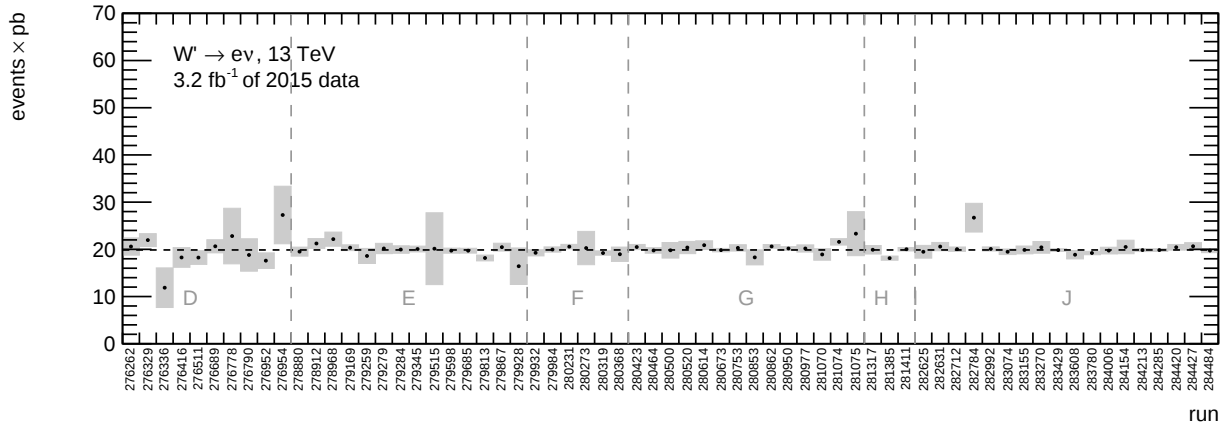
(a) $M_{W'} = 4 \text{ TeV}$



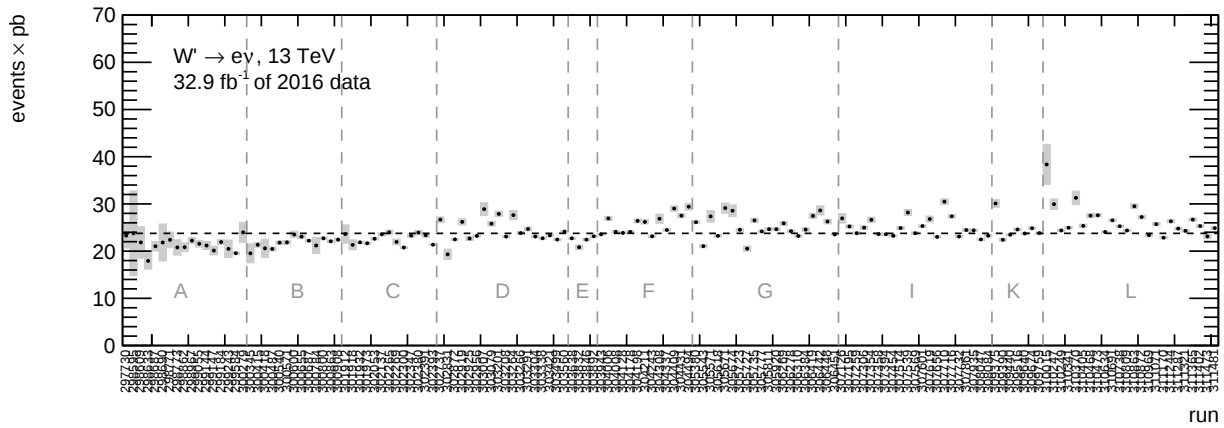
(b) $M_{W'} = 6 \text{ TeV}$

Figure E.1.: Signal templates for the electron channel, with interference effects, for various values of κ and two pole masses. The SM W contribution is included in the weight. The SM W estimate from a separate simulation is shown as well. The distributions are scaled to an arbitrary integrated luminosity of 1 fb^{-1} .

F. Event selection yield per run



(a) 2015



(b) 2016

Figure F.1.: Event yield per run in 2015 and 2016. The yield is defined as the number of events divided by the integrated luminosity of the run.

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