

# Interfering with the Higgs Boson

Lifetime Measurements and Effective Field Theory  
Interpretations of Higgs Boson Decays into Charged Vector  
Bosons at the LHC using the ATLAS Experiment



Bryan Alexander Kortman





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Experiment

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---

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## INTRODUCTION

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The Standard Model (SM) is *the* theoretical framework in particle physics. Developed to explain matter at its most fundamental level, the SM describes elementary particles and their interactions via three fundamental forces: the strong, weak, and electromagnetic forces. This model has been tested and verified across a vast array of experiments and energy scales, consistently showing excellent agreement with observed data. However, despite its successes, the SM falls short in explaining several key aspects of the universe such as dark matter, gravity, and the imbalance between matter and antimatter. These and more unresolved questions suggest the SM is still incomplete, sparking ongoing research aimed at the development of a more comprehensive theory of particle physics.

Historically, theoretical advancements have often led the way in particle physics, with theories predicting new subatomic particles and experiments designed to confirm these predictions. However, during the most recent runs of the Large Hadron Collider<sup>1</sup> (LHC), researchers did not find evidence for supersymmetry, a widely researched theoretical framework that many believed would be the next breakthrough after decades of development. As a result, the direction of current research has shifted. Instead of primarily searching for theoretical predictions, scientists are now focused on analysing data directly, in search of a deviation from SM predictions that might hint at new physics.

This thesis explores two promising avenues for uncovering new physics using  $pp$ -collision data at a centre-of-mass energy of  $\sqrt{s} = 13$  TeV, collected during Run-2 of the LHC. First, off-shell Higgs measurements offer a window into the rarely observed high-mass regime of the Higgs boson. This allows the measurement to estimate the Higgs lifetime and explore potential influences from undiscovered particles by directly measuring Higgs couplings. Since new physics models include extra scalars or a composite Higgs, the off-shell mechanism can deviate from the SM, altering the Higgs lifetime. Second, interpretations of Higgs measurements based on SM Effective Field Theory (SM-EFT) try to reveal possible higher-dimensional operators, which could point to BSM physics through the addition of a heavy new particle. Both methods leverage the quantum interference occurring in the production and decay of the Higgs boson and other SM particles. In these analyses, the observed interference patterns in final states are utilized to explore the potential existence of physics beyond the SM.

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### 1.1 Outline and personal contributions

All of the work presented in this thesis is conducted in close collaboration with the  $H \rightarrow WW^*$  analysis group, my ATLAS colleagues within the collaboration and the ATLAS group situated at the Nikhef institute in Amsterdam. This thesis aims to outline the main results of this work, focusing on my own contributions. Outlined below are the chapters in this thesis, the main projects I worked on during my PhD and my own personal contributions.

[Chapter 2](#) introduces the SM and describes its various particles and forces. After introducing the forces and particles, the chapter explains electroweak symmetry breaking and outlines the mechanism of the Higgs particle. It then briefly discusses the limitations of the SM and the need for a more complete theory.

[Chapter 3](#) discusses Higgs boson physics in the context of ATLAS measurements and the methodologies used in this thesis. It concludes with a detailed explanation of off-shell Higgs physics.

This chapter also details the study of off-shell Higgs boson samples for the off-shell  $H^* \rightarrow WW$  ATLAS analysis and the calculation of higher-order and massive quark loop corrections to the off-shell Higgs boson production cross-section. In collaboration with the authors of [1] and the LHC off-shell Higgs working group, I used an experimental setup to perform off-shell  $H^* \rightarrow WW$  predictions on the Nikhef stoombot cluster.

[Chapter 4](#) introduces the concept of SM-EFT and its application in ATLAS measurements to parameterize new physics effects.

[Chapter 5](#) provides an overview of the LHC and the CERN accelerator complex.

[Chapter 6](#) describes the ATLAS detector, detailing its sub-detectors and their specific roles in particle detection. This chapter also explains the trigger and data acquisition system used to store events for analysis.

#### Run-3 upgrades for L1 missing transverse momentum triggers

[Chapter 7](#) explains several upgrades to the ATLAS lowest level missing transverse momentum trigger and the new algorithms it supports. I developed and tuned L1 MET algorithms and implemented a trigger rate estimation model for L1 triggers. As part of the MET trigger group, I contributed to the performance assessment of new L1 MET trigger algorithms for Run-3, leading to the publication of expected performance plots for the ATLAS L1 calorimeter trigger at LHCP 2022 [2].

Collaboration with Radboud University on trigger studies led to the Green IT initiative, aiming to reduce the carbon footprint of experimental IT infrastructure [3].

[Chapter 8](#) describes the methods and algorithms used in the reconstruction of  $pp$ -collisions at the ATLAS experiment.

## **Analysis of off-shell Higgs boson production using $WW \rightarrow \ell\nu\ell\nu$ final states**

[Chapter 9](#) and [Chapter 10](#) detail the analysis of the off-shell Higgs boson using the  $H^* \rightarrow WW$  decay channel [4]. I kicked-off and led this analysis as one of two analysis contacts. I was directly involved in the development of the analysis strategy, implementing the (statistical) analysis within HWWAnalysisCode, estimation of uncertainties, estimation of the mis-identified lepton background from di-jet measurements, as well as other backgrounds.

## **Combined on-shell and off-shell $H \rightarrow WW$ measurement of the Higgs total width**

[Chapter 11](#) presents the combination of on-shell and off-shell Higgs boson measurements and interprets the results in terms of the Higgs boson width, lifetime and coupling strength ratios [4, 5]. In addition to the combination efforts, I contributed to the misidentified lepton and  $t\bar{t}$  background predictions for the on-shell analysis.

## **Combined SM-EFT interpretation of Higgs boson production and decay**

[Chapter 12](#) presents an SM-EFT interpretation of ATLAS Run-2 combined Higgs boson measurements which includes the  $H \rightarrow \gamma\gamma, ZZ^*, WW^*, \tau\tau$ , and  $b\bar{b}$  decay channels. I contributed to the interpretation of decay channel results in SM-EFT operators and defined the methodology for assessing SM-EFT acceptance effects, as published in [6].

## **Combined SM-EFT interpretation of Higgs and $WW$ measurements**

[Chapter 13](#) presents the first-ever Run-2 electroweak + Higgs combination and SM-EFT interpretation and combines two analyses with dissimilar techniques, as published in [7]. As an analysis contact, I initiated the analysis and contributed to its methodology, Monte Carlo simulation, and statistical implementation. This work underpins the combination of Higgs, electroweak, and electroweak precision observable measurements [8].

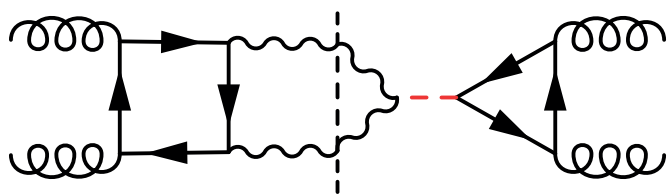
[Chapter 14](#) concludes and discusses the results presented in this thesis, followed by a technical summary and outlook in [Chapter 15](#).



PART I

THEORY





## THE STANDARD MODEL OF PARTICLE PHYSICS

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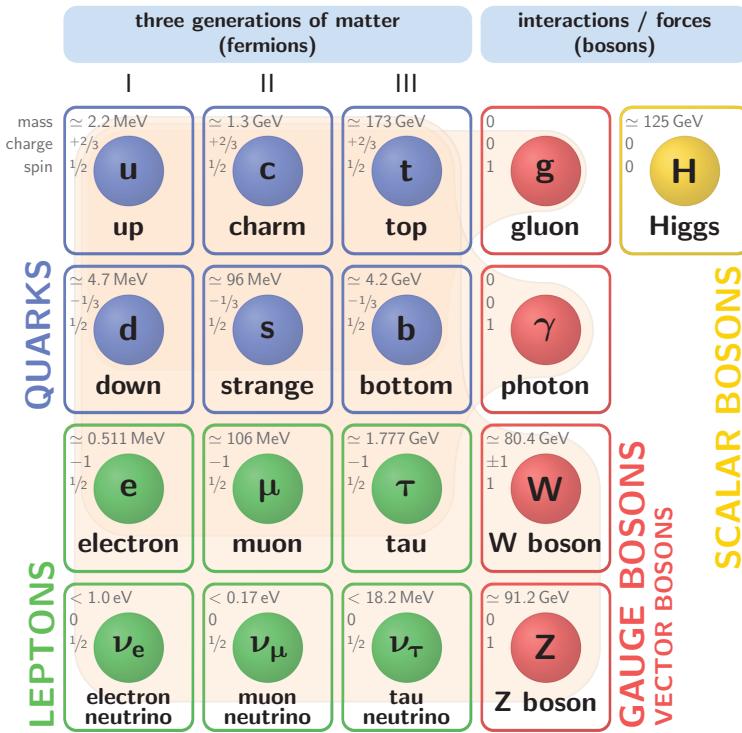
The SM of particle physics receives one of the most prominent positions in this thesis. accordingly, the following chapter describes the theoretical framework that models fundamental particles and predicts their interactions. It begins with an overview of the structure and particles included in the SM, followed by a detailed description of the different forces and the properties of the particles. The chapter concludes with a discussion on the limitations of the SM as a complete theory.

### 2.1 The contents of the SM

The SM of particle physics is a relativistic quantum field theory (QFT) and has been successful in describing the fundamental properties of matter and their development over time inside particle accelerator experiments. Since the discoveries of the first atoms and particles, various individual theoretical frameworks have been proposed and experimented with over the years in order to describe the kinematics of the different types of observed particles [9]. The combination of a set of these frameworks results in SM known today.

The first fundamental particle discovered is the electron by Joseph Thomson, using his cathode ray experiment in combination with a magnetic field [10]. This experiment enabled the calculation of mass to charge ratio ( $m/q$ ) of the cathode ray and the proposal that it consisted of small charged particles with a mass  $m_e$  and charge  $q_e = e$ . Many particle discoveries have followed since, along with composite particles containing more complex structure. However, since its discovery, the electron is still thought to be a structureless point-like particle and one of the elementary particles of nature.

In the next sections all fundamental particles featuring in the SM are introduced based on [9, 11]. These particles and their properties are illustrated in Fig. 2.1 and are subdivided into two groups, based on a quantum mechanical property, namely the intrinsic angular momentum, or spin, of the particles. Particles with a non-integer spin follow the Fermi-Dirac statistics and are called fermions; these particles cannot occupy the same quantum state, which means they can be distinguished from each other. The other group of particles has an integer spin and follows Bose-Einstein statistics; particles that carry integer spin are called bosons and are able to occupy the same quantum state which makes them indistinguishable.



**Figure 2.1:** A figure showing the fundamental particles of the SM. The leptons are shown in green, quarks in purple, the gauge bosons in red and the Higgs boson in yellow. The mass, spin and electromagnetic charge of each particle are shown as well. Taken from Ref. [12].

The interactions between the fundamental particles are mediated through the fundamental forces of nature; the dynamics incorporated in the SM describe three of the four fundamental forces. The first two are the strong and weak nuclear forces, the third force is the electromagnetic force, which allows electromagnetically charged particles to interact. The strong, weak and electromagnetic force are mediated via bosons with an integer spin; in the SM, these are named gauge bosons and are the gluon ( $g$ ), the charged vector boson ( $W^\pm$ ), the neutral vector boson ( $Z$ ) and the photon ( $\gamma$ ). The gluon is massless and mediates the strong force via the colour-charge it carries; the  $W^\pm$  and  $Z$  bosons mediate the weak force and are massive, significantly suppressing the interactions of the weak nuclear force; the electromagnetic force is mediated by the massless photon which carries no electromagnetic charge.

The fermion family consists of three generations of particles, each containing two quarks and two leptons, all with different masses and charges. The first generation contains the lightest particles making up the stable matter of our universe. The second and third generations structure is similar to the first generation except for their higher mass, which makes them unstable. The more massive fermion counterparts eventually decay into the first generation fermions.

The fermions are categorized into two groups based on their properties; the first group consists of quarks which interact with the strong force. The up-type quark carries an electromagnetic charge of  $+2/3e$  and comes in three generations or flavours; the up ( $u$ ), charm ( $c$ ) and top ( $t$ ) quarks. The down-type quarks carries an electromagnetic charge of  $-1/3e$  and also come in three generations, namely the down ( $d$ ), strange ( $s$ ) and bottom ( $b$ ) quarks. The leptons, carrying no colour-charge, are the already introduced electron ( $e$ ) along with the muon ( $\mu$ ) and tau ( $\tau$ ) particles; all leptons carry an electromagnetic charge of  $e$  in the SM. Each of the leptons can be produced in association with an electromagnetically neutral and similarly named neutrino; the neutrinos present in the SM are the electron neutrino ( $\nu_e$ ), the muon neutrino ( $\nu_\mu$ ) and the  $\tau$ -neutrino ( $\nu_\tau$ ).

The Higgs boson ( $H$ ) is the most recent addition to the SM following its discovery in 2012 at the LHC by the ATLAS and CMS experiments [13, 14]. It is not associated with a fundamental force of nature, but rather a particle generated through a mechanism responsible for giving mass to the fermions in the SM. In the SM, it is the only massive scalar particle, it carries no spin or charge and couples with all massive particles.

The SM uses lagrangian mechanics to describe the relativistic quantum fields that represent the fundamental particles in the theory. Comparable to non-relativistic lagrangian mechanics, the theory uses a relativistic lagrangian density  $\mathcal{L}(\psi_i, \partial_\mu \psi_i)$  to describe the equation of motion of the particles, fields and their interactions. The lagrangian density represents the summed kinetic and potential energy densities of the system, expressed in terms of a field  $\psi_i$  and its derivatives. By using the principle of least action, the equations of motion are determined by the Euler-Lagrange equations [11]:

$$\partial_\mu \left( \frac{\partial \mathcal{L}}{\partial (\partial_\mu \psi_i)} \right) - \frac{\partial \mathcal{L}}{\partial \psi_i} = 0.$$

This equation forms the basis for the derivation of the fermion field.

In the context of the SM, symmetries play a large role; It is found by the mathematician Emmy Noether that each continuous symmetry gives rise to a conserved Noether current and therefore a conserved charge, as stated by Noether's theorem [15]. A theory is symmetric under a group if its lagrangian is left invariant when a transformation belonging to that group is applied.

The fully realized SM is a non-commutative gauge field theory, based on the  $SU(3) \otimes SU(2) \otimes U(1)$  symmetry groups [16]. The  $SU(3) \otimes SU(2) \otimes U(1)$  symmetry groups have  $8 + 3 + 1 = 12$  generators with commutator algebra, where each of the generators constitute a transformation acting on the fields of the theory. The  $SU(2) \otimes U(1)$  group describes the electroweak interactions and the electric charge denoted by the quantum number  $Q$ . The  $SU(3)$  group is the colour group of the theory and is involved in strong interactions described in quantum chromodynamics (QCD).

## 2.2 Quantum electrodynamics

The Dirac lagrangian  $\mathcal{L}_{\text{Dirac}}$  forms the basis of quantum electrodynamics (QED); it describes the dynamics of massive spin-1/2 particles [16] and is given by

$$\mathcal{L}_{\text{Dirac}} = \bar{\psi} i \gamma^\mu \partial_\mu \psi - m \bar{\psi} \psi, \quad \bar{\psi} = \psi^\dagger \gamma^0.$$

In this relativistic lagrangian, the field  $\psi$  is the Dirac spinor, and  $\gamma^\mu$  represents the Dirac gamma matrices. The Dirac equation is the relativistic generalization of the Schrödinger equation and describes the dynamics of spin-1/2 fermions. It has four unique solutions, two of which are positive energy solutions and two are negative energy solutions. The negative energy solutions are interpreted as the anti-particles of the positive energy solutions. The existence of the anti-electron, or positron, was one of the first successes of the Dirac equation and was experimentally verified with the discovery of the positron in cosmic showers [17].

Because the Dirac lagrangian is invariant under the global phase transformation  $U(1)_{EM}$ , its fields and their derivatives transform as  $\psi \rightarrow e^{iq\alpha} \psi$  and  $\partial_\mu \psi \rightarrow e^{iq\alpha} \partial_\mu \psi$ . This corresponds to an overall phase change of the field  $\alpha$  and constant  $q$ , under which the Dirac lagrangian is invariant. However, when the transformation is made local, by replacing the phase change  $\alpha$  with a function of space-time  $\alpha(x^\mu)$ , the lagrangian is no longer invariant and the symmetry is broken.

$$\begin{aligned} \mathcal{L}'_{\text{Dirac}} &= e^{-iq\alpha(x)} \bar{\psi} i \gamma^\mu (e^{iq\alpha(x)} \partial_\mu \psi + e^{iq\alpha(x)} i q (\partial_\mu \alpha(x)) \psi) \\ &\quad - e^{iq\alpha(x)} e^{-iq\alpha(x)} m \bar{\psi} \psi \\ &= \mathcal{L}_{\text{Dirac}} - \bar{\psi} \gamma^\mu q (\partial_\mu \alpha(x)) \psi. \end{aligned}$$

In experiments, it is assumed that the physics model remains identical when applying a translation, making this symmetry a prerequisite for a proper theory of electrodynamics. Therefore, to restore the symmetry, an electromagnetic potential  $A_\mu$  is introduced as a field. The electromagnetic potential is a four-vector field and is introduced in the lagrangian as a covariant derivative defined as

$$\partial_\mu \rightarrow D_\mu \equiv \partial_\mu + iq A_\mu(x),$$

which contains the new vector field  $A_\mu$ . This field transforms as

$$A_\mu \rightarrow A_\mu(x) - \frac{1}{q} \partial_\mu \alpha(x)$$

when performing a local phase translation. By inserting the expression for  $A_\mu$  in the covariant derivative, the covariant derivative now transforms together with the local phase transformation of the field  $\psi$ :

$$\begin{aligned} D_\mu \psi &\rightarrow e^{iq\alpha(x)} (\partial_\mu \psi + i \partial_\mu \alpha(x) \psi + iq A_\mu \psi - iq \frac{1}{q} \partial_\mu \alpha(x) \psi) \\ &= e^{iq\alpha(x)} D_\mu \psi. \end{aligned}$$

As a consequence of the covariant derivative definition, the  $\psi^* D_\mu \psi$  terms, like the  $\psi^* \partial_\mu \psi$  terms under local transformation, are now phase invariant. With this substitution, the Dirac lagrangian, and any other real lagrangians, can be constructed while satisfying local phase symmetry. On top of this, since the lagrangian is required to be real, the field  $A^\mu$  must be physical as well [18]. The constant  $q$  is now identified as the charge of the fermion field and the gauge field  $A^\mu$  as the electromagnetic vector potential; The field  $A^\mu$  satisfies its own free lagrangian. The corresponding kinetic term for the electromagnetic vector field is

$$\mathcal{L}_{EM}^{\text{free}} = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu},$$

where  $F^{\mu\nu}$  is the field strength tensor,

$$F^{\mu\nu} = \partial^\mu A^\nu - \partial^\nu A^\mu = \begin{pmatrix} 0 & -E_x & -E_y & -E_z \\ E_x & 0 & -B_z & B_y \\ E_y & B_z & 0 & -B_x \\ E_z & -B_y & B_x & 0 \end{pmatrix}.$$

The vector field  $A^\mu$  is associated with the photon field and a mass term proportional to  $m_\gamma^2 A_\mu A^\mu$ . Taking a local gauge transformation of this term shows that it is only invariant if  $m_\gamma$  is equal to zero:

$$m_\gamma^2 A_\mu A^\mu \rightarrow m_\gamma^2 (A^\mu - \partial^\mu \alpha(x))(A_\mu - \partial_\mu \alpha(x)) \neq m_\gamma^2 A_\mu A^\mu.$$

Therefore, the requirement of local  $U(1)$  invariance automatically implies that the photon is massless. Finally, the theory of QED is described by the QED lagrangian

$$\mathcal{L}_{\text{QED}} = \bar{\psi} i \gamma^\mu \partial_\mu \psi - m \bar{\psi} \psi - q A_\mu \bar{\psi} \gamma^\mu \psi - \frac{1}{4} F_{\mu\nu} F^{\mu\nu}.$$

It describes the kinematics between a fermion, the electron with charge  $e$ , and the photon field. It is found that their interactions are proportional to the amplitude for a real electron to emit or absorb a real photon [18]. This proportionality is called the QED coupling constant, defined as  $\alpha_{QED} \equiv e^2 / 4\pi\epsilon_0 \hbar c \sim 1/137$ , where  $\epsilon_0$  is the permittivity of free space,  $\hbar$  is the reduced Planck constant and  $c$  is the speed of light [19].

### 2.3 Quantum chromodynamics

The strong force is mediated by the gluon field and is described by the QCD theory, which is based on the non-commutative  $SU(3)$  symmetry group. This symmetry group arises from the requirement that  $\mathcal{L}_{QCD}$  is invariant under the local  $SU(3)$  gauge transformation

$$\psi(x) \rightarrow \psi'(x) = e^{ig_s \alpha^a(x) T^a} \psi(x).$$

Here,  $T^a$  are the generators of the group, which are represented by the  $3 \times 3$  Gell-Mann matrices  $\lambda^a$  as  $T^a = 1/2 \lambda^a$ , along with the  $SU(3)$  gauge coupling  $g_s$  and  $\alpha^a(x)$ , the local gauge transformations corresponding to the eight generators [20, 21]. The covariant derivative for QCD can be written as

$$\partial_\mu \rightarrow D_\mu = \partial_\mu + ig_s G_\mu^a T^a,$$

where  $G_\mu^a$  must transform as

$$G_\mu^a \rightarrow G_\mu^a - \partial_\mu \alpha^a + i g_s G_\mu^a T^a.$$

Here, the last term arises from the non-abelian nature of QCD due to the generators not commuting. The  $SU(3)$  generators are related through their structure constants  $[T^a, T^b] \equiv T^a T^b - T^b T^a = i f_{abc} T^c$ . Using these ingredients, the lagrangian for a free quark field can now be written as

$$\mathcal{L}_{\text{quark}} = \bar{\Psi}(i\gamma^\mu \partial_\mu - m - g_s \gamma^\mu G_\mu^a T^a)\Psi.$$

In this lagrangian describing the quark dynamics, gauge field  $G_\mu^a$  is associated with the real gluon fields. Notice that  $\Psi$  has three components which correspond to the three colour states,  $\Psi = (\psi_1, \psi_2, \psi_3)$ , each of which is a Dirac spinor satisfying the free Dirac lagrangian. The QCD quark-gluon interaction vertex is identified in this lagrangian as

$$-g_s \bar{\Psi} \gamma^\mu G_\mu^a T^a \Psi.$$

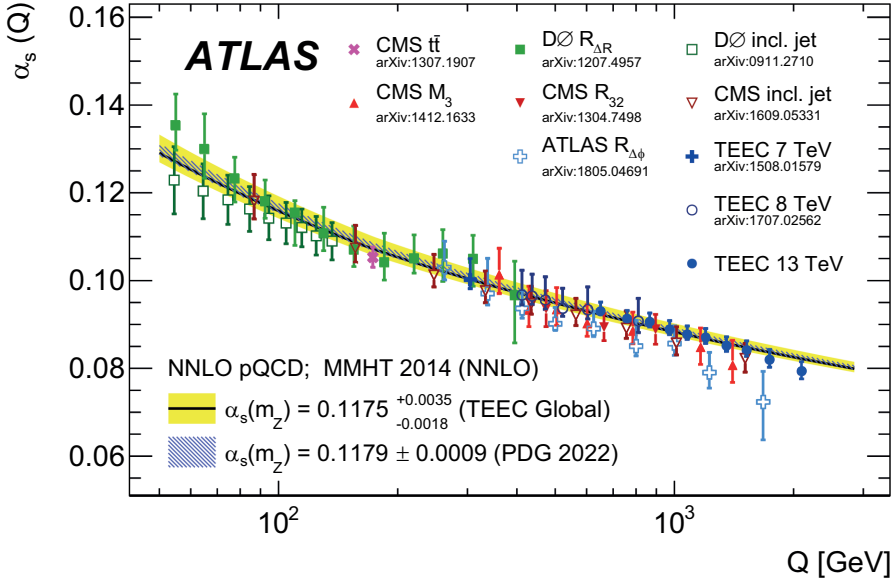
At last, the kinetic gauge-invariant term of the gauge bosons is added to the lagrangian, which is defined as

$$\mathcal{L}_{\text{kinetic}} = -\frac{1}{2} \text{Tr}[F^{\mu\nu} F_{\mu\nu}].$$

This kinetic term is proportional to the trace of the field tensors  $F_{\mu\nu} = \partial_\mu G_\nu - \partial_\nu G_\mu + i g_s [G_\mu, G_\nu]$ , with  $G_\mu = G_\mu^a T^a$  as the QCD gluon field strength tensor. In this definition, the non-vanishing structure constants  $f_{abc}$  appear, implying that gluons carry colour-charge and can therefore interact among themselves, giving rise to triple-gluon and four-gluon vertices. The  $SU(3)_C$  symmetry of QCD implies the existence of a conserved tri-valent colour-charge or Noether current, exchanged between quarks and eight gluons carrying the colour and anti-colour-charge in QCD vertices.

The QCD coupling constant  $\alpha_s$ , which is the amplitude of a quark interaction with a gluon, depends strongly on the momentum transfer  $q^\mu$  in the interaction. This can be shown in measurements of multijet events and is illustrated in Fig. 2.2. In this figure, the parameter  $Q = \sqrt{-q^2}$ , with  $q$  as the momentum transfer in the event, is the resolution energy used in the interaction and  $\alpha_s(Q)$  the coupling constant used in the simulation. The coupling constant  $\alpha_s$  becomes stronger as the momentum transfer in the event decreases. This asymptotic dependence on  $Q$  of  $\alpha_s$ , i.e. the fact that the strong (QCD) coupling becomes weak for processes involving large momentum transfers, leads to a feature of the chromodynamics theory called asymptotic freedom [22, 23].

A colour-charged particle, such as the quark, is continuously screened by virtual  $q\bar{q}$  and  $g\bar{g}$  matter-antimatter pairs that pop into and out of existence due to the asymptotic behavior of the QCD coupling. However, since the quark cloud retains a net colour charge, the QCD coupling constant and therefore the effective colour charge increases with distance as momentum transfer decreases.



**Figure 2.2:** Comparison of  $\alpha_s(Q)$  measurements with  $Q$  defined as the resolution of the event, proportional to the momentum transfer  $q$ . The value of  $Q$  is determined from fits to the transverse energy-energy correlations (TEEC) functions with the world average QCD prediction as input (hatched band) and the value obtained from the global fit (solid band) [24].

Ultimately, colour-neutral states are formed, made up from individual colour-charged carrying quarks, via the formation of hadrons. The effect, known as colour confinement [9], makes it impossible to isolate a colour-charged particle from neutral colour states.

Confinement makes it such that neutral colour states are the only states observed in nature; they can consist of a quark and an anti-quark, called mesons; or states of three quarks, called baryons. The process of hadronization, which forms hadrons out of quarks and gluons, can also produce more exotic colourless states, such as tetraquarks [25], containing two quarks and two anti-quarks, and pentaquarks [26], containing four quarks and one anti-quark.

## 2.4 Electroweak theory

In the SM, leptons are, in addition to their interactions with electromagnetically charged particles, able to interact with vector bosons. This facilitates weak interactions of leptons and caused the need to unify the weak and electromagnetic interactions between leptons and quarks in a unified electroweak theory.

The unified electroweak theory is formulated by Glashow-Weinberg-Salam [27–29], which has both  $U(1)$  and  $SU(2)$  invariance. The fermions in this relativistic theory are divided into three generations of left-handed ( $LH$ ), and right-handed ( $RH$ ), chiral quarks and leptons. For any Dirac field  $\psi$  the  $LH$  and  $RH$  chiral projections are defined



as

$$\begin{aligned}\psi_L &\equiv \frac{1}{2}(1 - \gamma^5)\psi, \\ \psi_R &\equiv \frac{1}{2}(1 + \gamma^5)\psi.\end{aligned}$$

In this equation, the Dirac matrix  $\gamma^5$  can be expanded as  $\gamma^5 = i\gamma^0\gamma^1\gamma^2\gamma^3$ . Dirac fields satisfying  $E \gg m$  correspond to the negative and positive helicity states respectively, defined by the projection of the spin of the particle onto its direction of motion. In electroweak theory, the  $RH$  fermion fields of each lepton and quark family are grouped into singlets and the  $LH$  into  $SU(2)$  doublets of Dirac spinors, while the neutrinos are assumed to be massless and only exist in their  $LH$  versions. This results in the following configuration of the fermion fields:

$$\begin{aligned}L_1 &= \begin{pmatrix} v_e \\ e^- \end{pmatrix}_L, & e_{R1} &= e_R^-, & Q_1 &= \begin{pmatrix} u \\ d \end{pmatrix}_L, & u_{R1} &= u_R, & d_{R1} &= d_R, \\ L_2 &= \begin{pmatrix} v_\mu \\ \mu^- \end{pmatrix}_L, & e_{R2} &= \mu_R^-, & Q_2 &= \begin{pmatrix} c \\ s \end{pmatrix}_L, & u_{R2} &= c_R, & d_{R2} &= s_R, \\ L_3 &= \begin{pmatrix} v_\tau \\ \tau^- \end{pmatrix}_L, & e_{R3} &= \tau_R^-, & Q_3 &= \begin{pmatrix} t \\ b \end{pmatrix}_L, & u_{R3} &= t_R, & d_{R3} &= b_R.\end{aligned}$$

In the electroweak theory, a  $SU(2)_L \otimes U(1)_Y$  symmetry is required. Under this symmetry, a  $LH$   $SU(2)$  doublet transforms as

$$\Psi \rightarrow \Psi'_L = e^{i\alpha(x)\mathbf{T} + i\beta(x)Y} \Psi_L.$$

In this equation,  $\mathbf{T} = \boldsymbol{\tau}/2$  are the  $SU(2)$  generators, where  $\boldsymbol{\tau}$  correspond to the Pauli matrices and  $Y$  to the generator for the  $U(1)$  symmetry. With this configuration, the  $RH$  components of the fields in the doublet transform only under the hypercharge  $Y$ :

$$\Psi \rightarrow \Psi'_R = e^{i\beta(x)Y} \Psi_R.$$

In this theory, the electric charge, derived in § 2.2, is connected with the sum of the third component of the weak isospin,  $I_3$ , and hypercharge  $Y$  of the fields:  $Q = I_3 + Y/2$ . The quantum numbers of the  $LH$  and  $RH$  leptons and quarks are given in Tab. 2.1. Like

**Table 2.1:** The electric charge  $Q$ , the isospin  $I_3$ , and the hypercharge  $Y$  for the  $LH$  and  $RH$  leptons and quarks.

|       | $v_L$          | $e_L$          | $e_R$ | $u_L$          | $d_L$          | $u_R$          | $d_R$          |
|-------|----------------|----------------|-------|----------------|----------------|----------------|----------------|
| $I_3$ | $+\frac{1}{2}$ | $-\frac{1}{2}$ | 0     | $+\frac{1}{2}$ | $-\frac{1}{2}$ | 0              | 0              |
| $Y$   | -1             | -1             | -2    | $+\frac{1}{3}$ | $+\frac{1}{3}$ | $+\frac{4}{3}$ | $-\frac{2}{3}$ |
| $Q$   | 0              | -1             | -1    | $+\frac{2}{3}$ | $-\frac{1}{3}$ | $+\frac{2}{3}$ | $-\frac{1}{3}$ |

the QED and QCD theory, the electroweak theory is based on the requirement that the lagrangian is invariant under local gauge transformations. The covariant derivative is defined as

$$D_\mu = \partial_\mu - ig \frac{\boldsymbol{\tau}}{2} \cdot \mathbf{W}_\mu - ig' \frac{Y}{2} B_\mu.$$

This results in the following lagrangian for the electroweak theory

$$\mathcal{L}_{\text{electroweak}} = \mathcal{L}_{\text{free}} - g \mathbf{W}_T^\mu \cdot \mathbf{b}_\mu - \frac{g'}{2} B_Y^\mu a^\mu.$$

In the  $\mathcal{L}_{\text{electroweak}}$  definition, the coupling constants  $g$  and  $g'$  belong to the  $SU(2)_L$  and  $U(1)_Y$  symmetry groups [30], respectively. The fields  $W_\mu$  and  $B_\mu$  are the gauge fields of the  $SU(2)_L$  and  $U(1)_Y$  symmetry groups. The  $W_\mu$  fields are the weak isospin gauge fields; the  $B_\mu$  fields are the weak hypercharge gauge fields that only interact with  $LH$   $SU(2)$  doublets that carry weak isospin and do not interact with  $RH$  singlets. The weak isospin gauge fields are the  $W_\mu^1$ ,  $W_\mu^2$  and  $W_\mu^3$  fields, which follow the structure of the three Pauli spin matrices  $\tau^1, \tau^2, \tau^3$ . Because of the non-diagonal elements of the  $SU(2)$  generators  $\tau^1, \tau^2$ , the physical weak charged current interaction fields are defined as

$$W^\pm \equiv \frac{1}{\sqrt{2}}(W_\mu^1 \mp iW_\mu^2).$$

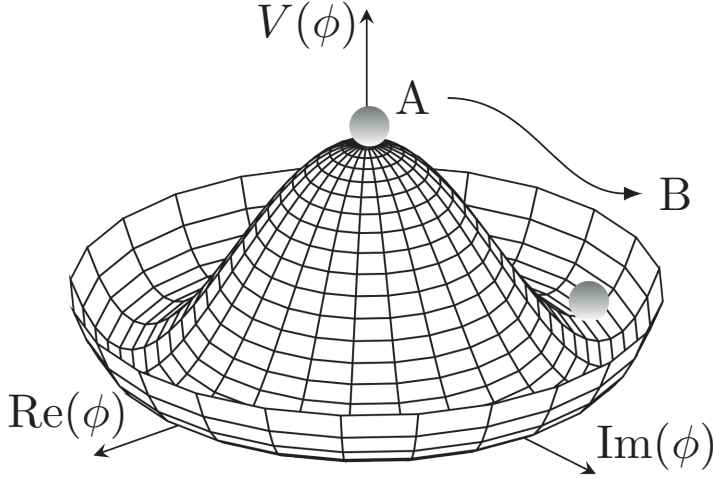
This leaves the  $W_\mu^3$  and  $B_\mu$  fields; together, these fields form the weak neutral current interaction fields, called  $Z$  and  $\gamma$ , via the weak mixing angle  $\theta_W$ . The weak mixing angle determines how the weak gauge fields mix to create the two neutral currents and is a fundamental parameter of the electroweak theory; it is defined as  $\cos(\theta_W) = g/\sqrt{g^2 + g'^2}$  and  $\sin(\theta_W) = g'/\sqrt{g^2 + g'^2}$ . Using these relations, the gauge fields  $A_\mu$  and  $Z_\mu$  are defined as

$$\begin{aligned} A_\mu &= B_\mu \cos(\theta_W) + W_\mu^3 \sin(\theta_W) \\ Z_\mu &= -B_\mu \sin(\theta_W) + W_\mu^3 \cos(\theta_W), \end{aligned}$$

The resulting photon field,  $A^\mu$ , was already discovered; the physical  $W^\pm$  and  $Z$  bosons were first observed at the UA1 and UA2 experiments [31] located at CERN. It was also experimentally verified that the  $W^\pm$  bosons only couples to  $LH$  particles. With this information, the first part of the SM lagrangian concerning the fermions is formulated to be:

$$\begin{aligned} \mathcal{L}_F &= -\frac{1}{4} W^{\mu\nu} W_{\mu\nu} - \frac{1}{4} B^{\mu\nu} B_{\mu\nu} \\ &+ \bar{L}_i i D_\mu \gamma^\mu L_i + \bar{e}_{Ri} i D_\mu \gamma^\mu e_{Ri} \\ &+ \bar{Q}_i i D_\mu \gamma^\mu Q_i + \bar{u}_{Ri} i D_\mu \gamma^\mu u_{Ri} + \bar{d}_{Ri} i D_\mu \gamma^\mu d_{Ri}. \end{aligned}$$

This lagrangian is invariant under the local  $SU(2)_L \otimes U(1)_Y$  transformation. The mass terms of the weakly interacting vector bosons are added by breaking the electroweak symmetry by the introduction of the Brout-Englert-Higgs mechanism [32–34]. The fermionic mass terms are introduced with the help of gauge-invariant Yukawa interactions of the fermions with the Higgs field.



**Figure 2.3:** The 3D Mexican hat potential,  $V(\Phi) = \mu^2|\Phi|^2 + \lambda|\Phi|^4$ , illustrating the spontaneous symmetry breaking of the system while transitioning from state A to state B. Taken from Ref. [12].

## 2.5 The Brout-Englert-Higgs mechanism

In this section, the procedure of spontaneous symmetry breaking of the electroweak theory via the Higgs mechanism is introduced [35]. By adding a scalar Higgs field to the theory, the gauge bosons acquire mass through their interactions. The newly added field is a weak isospin doublet of two complex scalar fields:

$$\Phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} \phi_1 + i\phi_2 \\ \phi_3 + i\phi_4 \end{pmatrix}.$$

The field  $\Phi$  is a  $LH$  doublet similar to the electron-neutrino doublet. It has a weak isospin  $1/2$  and the charges of the upper and lower component are chosen such that the hypercharge  $Y = 1$ . The potential  $V(\Phi)$ ,

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

is chosen with  $\mu^2 < 0$ , which is added to the lagrangian. The Mexican hat potential, shown in Fig. 2.3, receives its name from the fact that it looks like a smooth Mexican hat centred around the  $\hat{z}$ -axis, symmetric with respect to its centre  $\hat{z}$ -axis. The top origin of the potential is an unstable minimum for the system and it may spontaneously fall to either side of the potential, towards the point of lowest energy. With this transition, the total system loses its  $\hat{z}$ -axis symmetry in state A and transitions to state B. This transition is called spontaneous symmetry breaking and is the mechanism that gives mass to the weak gauge bosons [36].

Using the new isospin doublet and Mexican hat potential, the lagrangian is written as

$$\mathcal{L} = (D^\mu\Phi)^\dagger(D_\mu\Phi) - V(\Phi),$$

where  $D_\mu$  is the covariant derivative associated to the  $SU(2)_L \otimes U(1)_Y$  symmetry, previously defined in § 2.4. Since the shape of the underlying potential depends on the values of  $\mu$  and  $\lambda$ , the limit of  $\lambda > 0$  is required to ensure the potential has a low bound, in order to ensure a positive non-zero energy state. In case of a positive  $\mu^2$ , the potential has a single minimum at  $\phi_0$ . When  $\mu^2 < 0$ , there is no single minimum in the potential energy located at  $(0, 0)$ . Instead, a continuous minimum is present in the circle of points satisfying  $\sqrt{Re\Phi^2 + Im\Phi^2} = \sqrt{-\mu^2/\lambda} = v$ .

Now the lagrangian is defined, a choice for a vacuum is made without losing generality. The fields  $\phi_1 = \phi_2 = \phi_4 = 0$  and  $\phi_3 = v + H$  are chosen such that  $H$  represents the actual Higgs field. The vacuum expectation value of the Higgs field is defined as

$$\text{Vacuum} = \phi_0 = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + H \end{pmatrix},$$

with

$$\phi_0^\dagger \phi_0 = \frac{1}{2}(v^2 + H^2) = -\frac{\mu^2}{\lambda}.$$

When the Higgs field is expanded around zero, which introduces the Goldstone bosons  $\phi_1, \phi_2, \phi_4$ , the field is written as

$$\Phi(x) = \frac{1}{\sqrt{2}} \begin{pmatrix} \phi_1 + i\phi_2 \\ \frac{1}{\sqrt{2}}(v + H) - i\phi_4 \end{pmatrix} \approx e^{i\theta_a(x)\tau^a(x)/v} \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ \frac{1}{\sqrt{2}}(v + H(x)) \end{pmatrix}$$

The Goldstone bosons  $\phi(x)_{1,2,4}$  are non-physical states and arise when a global symmetry is spontaneously broken, stated in the Nambu-Goldstone theorem [37, 38]. These bosons are gauged away by choosing the transformation of the field accordingly and are then absorbed into an overall phase factor. This is called the unitary gauge  $\phi_i(x) = 0$  which makes the remaining Higgs interactions become longitudinal components of the weak gauge bosons.

This mechanism leads to the generation of the mass terms of the gauge bosons, which are obtained taking the electroweak covariant derivative of the field  $\phi$  [30]:

$$\begin{aligned} |D_\mu \phi|^2 &= \frac{1}{2} \left| \left( \partial_\mu - ig \frac{\boldsymbol{\tau}}{2} \cdot \mathbf{W}_\mu - ig' \frac{B_\mu}{2} \right) \begin{pmatrix} 0 \\ v + H \end{pmatrix} \right|^2 \\ &= \frac{1}{2} (\partial_\mu H)^2 + \frac{1}{8} g^2 (v + H)^2 |W_\mu^1 + W_\mu^2|^2 + \frac{1}{8} (v + H)^2 |gW_\mu^3 - g'B_\mu|^2. \end{aligned}$$

In this equation,  $\mathbf{W}_\mu$  is first expanded in terms of its weak isospin gauge fields; after which the equation is rewritten into  $W^\pm$ , in order to represent the physical fields and reveal the generated mass term. The expected longitudinal mass term of a charged gauge boson appears as  $M_W^2 W_\mu^+ W^{-\mu}$ , and the mass is identified to be

$$M_W = \frac{1}{2}vg.$$

The remaining term is comprised of the fields  $W_\mu^3$  and  $B_\mu$ , from which after substitution of the new fields

$$A_\mu = \frac{g'W_\mu^3 + gB_\mu}{\sqrt{g^2 + g'^2}}, \text{ with } M_A = 0,$$

$$Z_\mu = \frac{gW_\mu^3 - g'B_\mu}{\sqrt{g^2 + g'^2}}, \text{ with } M_Z = \frac{1}{2}v\sqrt{g^2 + g'^2},$$

the last two neutral gauge boson mass terms can be identified. This results in the following three bilinear mass terms generated by the spontaneous  $SU(2)_L \times U(1)_Y \rightarrow U(1)_Q$  symmetry breaking:

$$M_W^2 W_\mu^+ W^{-\mu} + \frac{1}{2}M_Z^2 Z_\mu Z^\mu + \frac{1}{2}M_A^2 A_\mu A^\mu.$$

Now the weak mixing angle given earlier is reintroduced as

$$\frac{g'}{g} = \tan \theta_W,$$

with  $g$  and  $g'$  the coupling strength of the  $SU(2)_L$  and  $U(1)_Y$  interaction respectively. The spontaneous breaking of the  $SU(2)_L \otimes U(1)_Y$  symmetry is complete and the  $W^\pm$  and  $Z$  bosons acquire masses. The  $U(1)_Q$  symmetry remains unbroken and thus its generator, the photon, remains massless. Inserting the mass terms and working out the full lagrangian generates interactions with the Higgs boson. At last, the Higgs boson lagrangian is expressed in real particle masses, coupling parameters and the newly defined weak mixing angle:

$$\begin{aligned} \mathcal{L}_{\text{Higgs}} = & \frac{1}{2}\partial_\mu H^\mu \partial^\mu H_\mu + M_H^2 H_\mu H^\mu \text{ (Higgs kinetic and mass term)} \\ & + M_W^2 W_\mu^+ W^{-\mu} + M_Z^2 Z_\mu Z^\mu \text{ (W and Z mass terms)} \\ & + gM_W H W_\mu^+ W^{-\mu} + \frac{g^2}{4} H_\mu H^\mu W_\mu^+ W^{-\mu} \text{ (HWW and HHWW coupling)} \\ & + \frac{gM_Z}{2\cos(\theta_W)} H Z_\mu Z^\mu + \frac{g^2}{4\cos(\theta_W)^2} H_\mu H^\mu Z_\mu Z^\mu \text{ (HZZ and HHZZ coupling)} \\ & - \frac{gM_H^2}{4M_W} H^3 - \frac{g^2 M_H^2}{32M_W^2} H^4 \text{ (HHH and HHHH coupling)}. \end{aligned}$$

The physical real field  $H$  describes the Higgs boson, a neutral scalar particle with has a mass of  $M_H = \mu/\sqrt{2}$ . Its triple and quartic self interactions are proportional to  $M_H^2$ . In this lagrangian, the tri-linear  $HW^+W^-$ ,  $HZZ$  and quartic  $HHW^+W^-$ ,  $HHZZ$  couplings are also incorporated. The Higgs coupling to the gauge bosons is identified to be directly proportional to the gauge boson mass

$$g_{HWW} = gM_W, \quad g_{HZZ} = \frac{g}{\cos(\theta_W)} M_Z.$$

## 2.6 Yukawa terms and fermion masses

The complex doublet field  $\Phi$ , used to generate the masses of the gauge bosons, is similarly used to generate fermion masses: mass terms of the form  $\bar{\psi}m\psi = m(\bar{\psi}_L\psi_R + \bar{\psi}_R\psi_L)$  become part of the electroweak lagrangian discussed in § 2.4 due to this mechanism. Together with its charge conjugated partner,  $\Phi^C = 1/\sqrt{2} \begin{pmatrix} v + H & 0 \end{pmatrix}$ , the doublet field generates masses which take the form of Yukawa interactions [39] after the spontaneous symmetry breaking. The Yukawa terms are added to the lagrangian and are defined as

$$\mathcal{L}_{\text{Yukawa}} = -\lambda_{ij}\bar{e}_{Li}\Phi e_{Rj} - \lambda'_{ij}\bar{Q}_i\Phi u_{Rj} - \lambda''_{ij}\bar{Q}_i\Phi d_{Rj} + h.c.$$

In this equation, the couplings  $\lambda_{ij}$ ,  $\lambda'_{ij}$  and  $\lambda''_{ij}$  are the Yukawa coupling constants and the indices  $i$  and  $j$  represent the three generations of quarks and leptons. The terms labelled  $L$ ,  $Q$  and  $e$ ,  $u$ ,  $d$  are the  $LH$  lepton and quark doublets and  $RH$  singlets, respectively. The Yukawa couplings of the up-type and down-type quarks are labelled  $\lambda'_{ij}$  and  $\lambda''_{ij}$  and take the form of  $3 \times 3$  matrices with off diagonal elements allowing for the mixing of the generations.

Because of the diagonal elements, there is a mismatch between mass eigenstates and weak eigenstates, allowing states to transition between them. The Cabibbo-Kobayashi-Maskawa (CKM) matrix for quarks [40, 41] is a unitary matrix that describes the mixing of the quark weak eigenstates with the mass eigenstates. The matrix can be parameterized by three angles and one phase, which is the only source of CP violation in the SM [11]. By diagonalizing the Yukawa couplings to the mass eigenstates, their coupling strength to the Higgs boson is obtained and therefore their mass. The mass of the fermions is then given by

$$m_f = \frac{\lambda_{ij}v}{\sqrt{2}}, \quad m_{u,c,t} = \frac{\lambda'_{ij}v}{\sqrt{2}}, \quad m_{d,s,b} = \frac{\lambda''_{ij}v}{\sqrt{2}}.$$

## 2.7 Theoretical limitations of the SM

The SM has been shown to describe fundamental particles and their interactions in controlled experiments with a high degree of precision. Moreover, it has predicted the existence of certain particles prior to their discovery. Although the SM is presented as a theory describing the kinematics of subatomic particles and the matter making up the our world, it has several limitations regarding the matter it seeks to describe.

Firstly, the SM does not account for gravity. The gravitational force acting on sub-atomic particles is not included in the SM and is currently only described by the theory of general relativity. Since the SM is a QFT and general relativity is a classical field theory, the two are not easily combined. This necessitates the development of a quantum theory of gravity to describe the universe at the smallest scales, which can then be integrated with the SM.

Efforts are made to develop a quantum theory of gravity, introducing a boson called the graviton, which mediates the gravitational force [42]. The graviton is predicted to be

a massless particle with a spin of 2. However, when studying quantum gravity at the Planck scale, current theories become unrenormalizable, predicting infinite values for certain quantities. Quantizing classical general relativity remains a significant challenge in theoretical physics [43].

Secondly, the relatively low mass of the Higgs boson (125 GeV) poses a problem. Considering the interactions of the Higgs boson, its mass receives corrections up to the Planck scale. Hence, it is expected that corrections from many high-mass BSM fermions would significantly increase its mass. Only through many higher-order cancellations can the Higgs boson acquire a mass close to the weak scale. This results in a highly tuned theory, suggesting that electroweak scale physics depends on the interaction structure of high-mass BSM fermions. This effect breaks the principle of separation of scales, which states that low-energy physics should behave independently of high-energy physics and is known as the hierarchy problem [44].

Another shortcoming of the SM is its inability to explain the origin of the matter-anti-matter asymmetry in the universe. Observations show that the universe is composed almost entirely of matter, a phenomenon known as baryon asymmetry [45]. The complex phases in the CKM and PMNS matrices are unlikely to cause enough CP violation in order to explain these observations.

Another discrepancy the SM fails to account for is the observation of neutrino oscillations [46], which imply that neutrinos have mass. The SM includes only massless neutrinos, directly contradicting observed neutrino oscillations. Neutrino oscillations can be incorporated into the SM by adding the Pontecorvo-Maki-Nakagawa-Sakata matrix [47, 48], similar to the CKM matrix, describing the mixing of neutrino weak eigenstates with mass eigenstates.

Two additional phenomena not explained by the SM are the existence of dark matter and dark energy [9]. The cosmological observations of these concepts lack explanation within the current SM. Dark Matter is suspected to constitute the majority of the mass of the universe through studies of galaxy rotation curves and gravitational lensing of galaxy clusters and is subject to various searches using accelerators or cosmic particles [49, 50]. Dark energy is thought to be the cause of the universe's accelerated expansion [9].

## 2.8 Phenomenology of particle physics at hadron colliders

In this thesis, the SM is used to describe the particle kinematics at the LHC. Hadron colliders, like the LHC, have the ability to reach high centre-of-mass energies and produce rare interactions or particles. However, because of the composite nature of hadrons, the event structure of a single collision is complex and typically involves multi-particle calculations. In order to describe how collisions are modelled at hadron colliders, the following sections break down the process of a collision event into its constituent parts.

### 2.8.1 From theory to experiment

The production of a particle with a high invariant mass or momentum transfer can be described by factorizing the different stages a collision undergoes. The interaction rate of processes at collider experiments is expressed in the form of a total cross-section and unit of area. The total cross-section can be viewed as a measure of probability for a given process to occur; in the next examples, the production of a Higgs boson at a proton-proton collider is used. The total cross-section is given by the convolution

$$\sigma_{p_1 p_2 \rightarrow H} = \sum_{a,b} \int_0^1 dx_a dx_b \int d\Phi_H f_a^{p_1}(x_a, Q^2) f_b^{p_2}(x_b, Q'^2) \\ \times \frac{A}{2\hat{s}} |\mathcal{M}_{pp \rightarrow H}|^2(\Phi_H, Q^2, Q'^2),$$

with,

$$\frac{A}{2\hat{s}} |\mathcal{M}_{pp \rightarrow H}|^2(\Phi_H, Q^2, Q'^2) = \frac{d\sigma_{p_1 p_2 \rightarrow H}(\Phi_H, Q^2, Q'^2)}{d\Phi_H}.$$

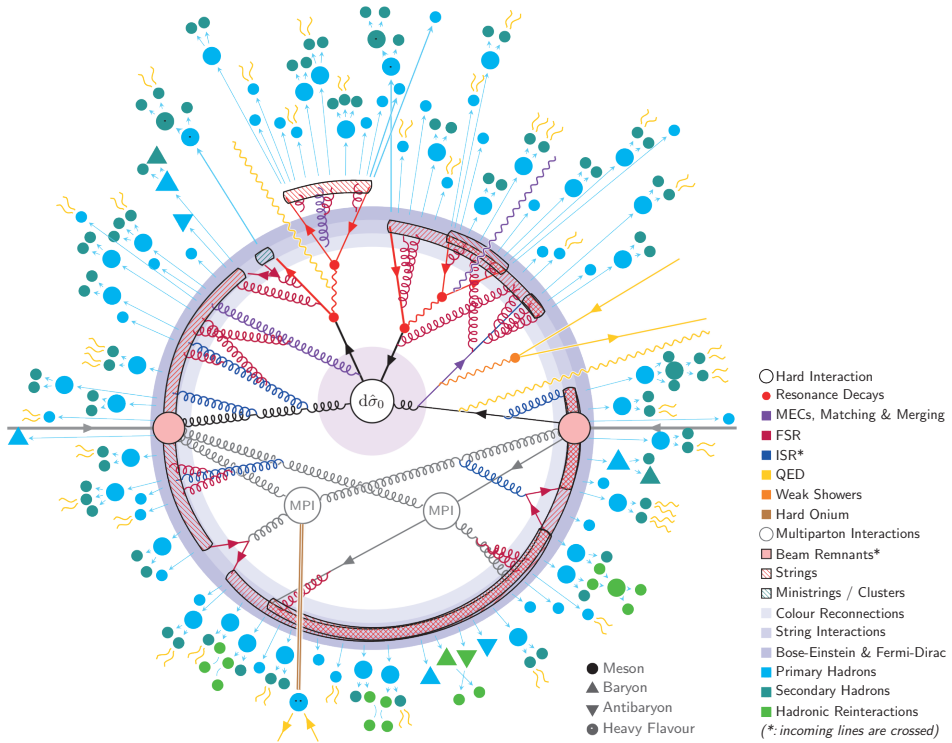
The functions  $f_a^{p_1}(x_a, Q^2)$  and  $f_b^{p_2}(x_b, Q'^2)$  in this equation are the parton distribution functions (PDFs) of the incoming protons [51–53]. The PDFs describe the probability of a parton  $a$  or  $b$  with a fraction  $x_a$  or  $x_b$  of the proton's momentum. The PDFs depend on the momentum transfer  $Q^2$  and  $Q'^2$ , and are determined by global fits to experimental data from deep inelastic scattering experiments [54]. The  $A/2\hat{s} \cdot |\mathcal{M}_{pp \rightarrow H}|$  term is the matrix element for the process of interest with a pre-factor  $A$  and  $\Phi_H$  is the phase-space of the final state  $H$ , which is the space of all possible final state momenta of  $H$  determined by the kinematics of the process. The total centre-of-mass energy in the rest frame is defined as  $\hat{s}$ . The total cross-section is now identified as the sum of all possible parton-parton interactions that give rise to a Higgs boson  $H$ , which then used to calculate the number of expected events in the detector.

### 2.8.2 Partons and matrix elements

The process of a  $pp$ -collision producing a  $t\bar{t}$  pair, one of the largest backgrounds in this thesis, is illustrated in Fig. 2.4. This process, which is characterized by a significant parton activity, consists of several parts. First, two partons inside the hadron interact and transfer momentum. Therefore, there is need for a detailed description of the accelerated protons and colliding hadrons, achieved by the PDFs. Secondly, the momentum transfer is described by the matrix element, named the hard-scattering process.

The hard-scattering process is followed by the production of a large number of soft and collinear partons, which are modelled using a parton shower. The parton shower is a semi-classical description of the radiation of partons from the hard-scattering and often requires tuning to the experiment of interest. The underlying event is also modelled in the collision, which consists of the soft and collinear radiation that is not part of the initial hard-scattering. This is followed by the hadronization of the colour-carrying partons.

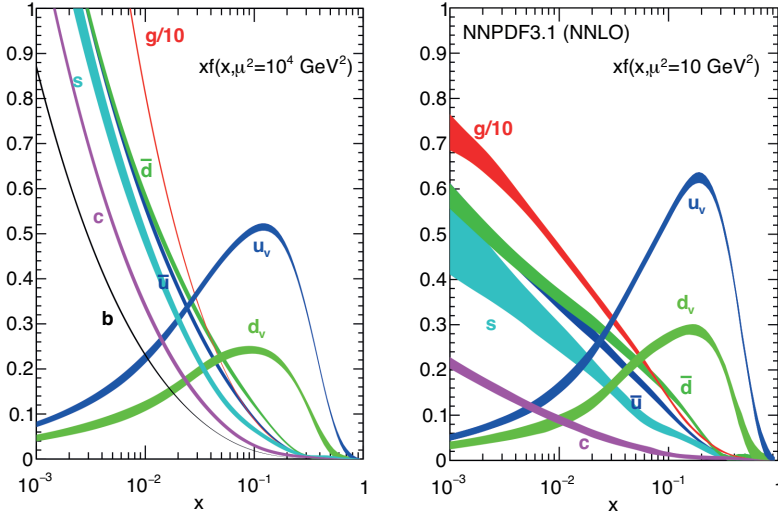




**Figure 2.4:** The figure illustrates the structure of a  $pp \rightarrow t\bar{t}$  event, which is a significant background in  $H^* \rightarrow WW$  analyses and simulated using Pythia8. The event structure is simplified for clarity. There are fewer shower branchings and final-state hadrons compared to actual Pythia8 events. Recoil effects are not depicted with complete accuracy. Additionally, weak decays of light-flavour hadrons are excluded. Lastly, the beam remnants and the incoming lines' radiation branchings, both before and after branching, should be interpreted with reversed momentum [55].

At last, the produced hadrons and leptons, which includes the underlying event, are detected in the detector. Since every proton is composed of three valence quarks surrounded by a sea of gluons and quark-anti-quark pairs, the PDFs representing the incoming protons are a function of the momentum fraction  $x$  of the proton carried by the colliding parton and the momentum transfer  $Q^2$ . In this distribution function, each parton in the proton carries a fraction  $x_i \in [0, 1]$  of the total proton momentum, with  $p_i = x_i p$  the momentum of the parton.

Since the content of a proton cannot be described by perturbation theory, due to the non-perturbative nature of the partons inside the proton; the PDFs and their dependence on  $x$  are determined by global fits of experimental data to the DGLAP equations [11]. Typically, at low  $Q^2$  the PDFs are dominated by the valence quarks and at high  $Q^2$  by gluons. High  $x$  is carried often by the valence quarks or, at high  $Q^2$ , sea quarks. In Fig. 2.5, the PDFs of the proton are shown at low and high momentum transfer  $Q^2$ , or scale  $\mu^2$ , illustrating the probabilities of colliding with a gluon, valance or sea quarks.



**Figure 2.5:** The PDFs of the proton. The left plot shows the PDFs of the proton at a low momentum transfer  $Q^2$  and the right plot shows the PDFs at a high momentum transfer  $Q^2$ . These PDFs are determined by global fits of experimental data to the DGLAP equations by the NNPDF collaboration [56].

The stage where new fields or particles are produced, the hard-scatter process modelled by the matrix element, must closely replicate the corresponding hard-scatter event observed in the detector. The PDFs  $f_a^{p_1}(x_a, Q^2)$  and  $f_b^{p_2}(x_b, Q'^2)$  supply the probability of finding a parton with momentum  $p_1$  and  $p_2$  at the central energy scale of  $Q^2$ . The factorization scale  $Q^2$  is the scale at which the parton kinematics are described, based upon the factorization theorem [57]. The scale at which the hard-scattering process takes place is  $\mu_R$  and is close to the scale of the momentum transfer  $Q$  of the hard-scatter. In hadron colliders, the QCD coupling  $\alpha_s(\mu_R \approx Q^2)$  gives an indication of the strength of the strong interaction in the process.

During the collision, real and virtual gluons are produced which give rise to collinear singularities [11]. These singularities arise when the momentum of a parton becomes collinear to the momentum of the proton, in the QCD theory they are absorbed into the renormalization scale  $\mu_R$ . This renormalization scale ensures infrared (IR) safety [58] and enables the calculation of IR-safe observables, which are insensitive to soft and collinear parton branchings and typically integrated over the full phase space. The total cross-section [59] is one of these IR-safe observables:

$$\sigma_{p_1 p_2 \rightarrow H} \propto \sum_{a,b} \int_0^1 dx_a dx_b f_a^{p_1}(x_a, \mu_F^2) f_b^{p_2}(x_b, \mu_F^2) \times \sigma_{ab \rightarrow H}^{\text{total}}(x_a, x_b, \mu_R^2).$$

Since the matrix element is computed perturbatively, the total cross-section  $\sigma_{ab \rightarrow H}^{\text{total}}$  includes the leading order (LO) contribution, which is the most significant, followed by higher-order corrections:

$$\sigma_{ab \rightarrow H}^{\text{total}} = \sigma_0 + a_1 \sigma_1 + a_2 \sigma_2 + \dots$$

In this equation,  $\sigma_0$  represents the LO contribution,  $\sigma_1$  denotes the next-to-leading-order (NLO) term, and  $\sigma_2$  indicates the next-to-next-to-leading-order (NNLO) term. The coefficients  $a_i$  scale these terms and can be either positive or negative. Higher-order corrections to the LO term are calculated using specialized programs and face two primary challenges [59, 60]. The first challenge is the factorial increase in the number of Feynman diagrams when accounting for colour structure, Lorentz structure, and all final state particles. The second challenge arises from the growth in the number of divergences in the integrands as more Feynman diagrams are included. Additionally, adding extra virtual and real particles in higher-order terms makes phase-space integrals more difficult and numerically intensive to evaluate. The choice of the renormalization scale  $\mu_R$  and the factorization scale  $\mu_F$  depends on the specific process and reflects the energy scale at which the process occurs [61].

Once the scale is chosen, uncertainties are assigned to account for the perturbative higher-order corrections. Including higher-order terms in the matrix element reduces the impact of the  $\mu_R$  choice, as it absorbs fewer higher-order perturbative corrections. Similarly, if the PDFs accurately describe the partons at higher-orders, the effect of  $\mu_F$  is minimized because the non-perturbative terms are smaller. Higher-order corrections, which are challenging to compute, are often estimated using an inclusive ( $K$ ) factor or a differential K-factor ( $K(x)$ ) [62]. The total cross-section can then be simplified to  $\sigma_{ab \rightarrow H}^{\text{total}} = \sigma_0 \cdot K(x)$ . This enables experimental particle physics to incorporate approximated higher-order corrections that exceed the current capabilities of the generator, often including high-scale uncertainties as an added consideration.

### 2.8.3 Parton showers, matching and merging

Parton showers model the approximate higher-order real-emission corrections to hard processes using the factorization properties of the perturbative matrix element. They estimate the virtual corrections to the process of interest and allow for the evaluation of observables sensitive to soft collinear radiation. The showering happens in an ordered fashion, splitting the particle with the highest momentum first and continuing to the next particle [61].

Classic parton showers use fragmentation models in order to model the hadronization of the partons into hadrons non-perturbatively. In order to further improve the description of certain observables, parton showers often go beyond these approximations. In a process called matching [63, 64], the parton shower expression is computed at fixed-order and is subtracted from the higher-order hard process calculation; this subtraction is

carried out to eliminate double counting of higher-order effects, and ensures that only the subtracted result is further processed by the parton shower.

In the method called merging [63, 64], a separate calculation is performed for each parton multiplicity of interest. By variation of a momentum cutoff selection on the partons, the soft and collinear divergences of the hard matrix elements are regulated. The parton shower is combined with all parton multiplicity calculations and potential double counted events are removed by vetoes on shower branchings by the fragmentation model.

A generated event is often accompanied with accelerated colour and electromagnetic charges, and thereby bremsstrahlung can occur [61]. Emissions that are associated with the incoming colliding partons are called initial-state radiation (ISR) and emissions that are associated with outgoing partons are instead called final-state radiation (FSR). Both are non-trivial to model and are approximated by the parton shower.

The goal of the techniques is that they provide a detailed description of the final state so that any experimental observable can be predicted and compared with data. On top of this, they are crucial for the development of a proper understanding of the signal and background. This, in turn, allows the two to be treated separately in analyses and enhances the overall sensitivity.

## 2.8.4 Monte Carlo generators and detector simulation

There are several Monte Carlo generators used throughout this thesis, each of them with their own capabilities. Monte Carlo programs can specialize in a certain type of process or final state. However, there are also general purpose generators which can simulate generic processes as long as they involve particles implemented in the theoretical model. All generators simulate the hard-scatter interaction first, after this simulation a parton shower is needed to simulate the full  $pp$ -collision. This section aims to outline the key generators and their most important features vital to this work.

MadGraph5\_aMC@NLO [65, 66], a parton-level event generator that simulates a hard-scattering process at LO and NLO in QCD. Additionally, NLO EW corrections can be added if the supplied model allows it.

Pythia8 [55, 67] is a general purpose event generator that simulates the full  $pp$ -collision including parton showering and hadronization. A model for the underlying event and multiple parton interactions is included.

Herwig7 [68] is a LO parton shower generator typically used to study an angular ordered parton shower, it also includes a model for the underlying event and multiple parton interactions.

Sherpa [69] is another general purpose event generator that simulates the full  $pp$ -collision including parton showering, underlying event, multiple parton interactions and hadronization.

2 The Powheg-Box and Powheg-Box-Res [70–73] are parton-level event generator frameworks that simulate a hard-scattering process at NLO in QCD, including parton shower. They can also be interfaced with a different parton shower and hadronization model when needed. The Powheg-Box-Res version of Powheg is specialized in the generation of radiation in decaying resonances, which can be challenging to model with general purpose generators.

After simulation of the particle dynamics, the response of the ATLAS detector is simulated with GEANT4 [74]. This program simulates the material response to the particles passing through, along with the subsequent readout of the detector. The output of the simulation is processed with the same software and analysis as the  $pp$ -collision data, which allows for comparison between the simulated and real data.

## HIGGS BOSON PHYSICS AND PHENOMENOLOGY

Because the Higgs boson takes centre stage in this thesis, this chapter provides an overview of the Higgs boson in the SM, its production and decay channels. After introducing the Higgs boson properties, the analysis methods used throughout this thesis are introduced. The chapter also discusses the importance of off-shell Higgs physics and the Higgs lifetime, which are one of the main subjects of this thesis. At last, the chapter introduces the modelling of the Higgs processes in the off-shell measurement.

### 3.1 The SM Higgs boson lifetime

Since the Higgs boson couplings depend on its mass and the Higgs boson mass is a free parameter of the SM, it is an important property to measure. The latest Higgs boson mass measurement, shown in Fig. 3.1, is performed with sub-per mille precision to be  $m_H = 125.11 \pm 0.11$  GeV [75]. From the measured free parameter  $m_H$ , the Higgs boson width and the SM coupling strengths can then be determined. The assumption of a 125 GeV Higgs mass results in a predicted Higgs width of  $\Gamma_H = 4.1 \pm 0.17$  MeV [76], this is a very narrow width compared to the mass of the Higgs boson.

This narrow resonance allows to apply the narrow width approximation and results in the factorization of the total Higgs cross-section by replacing the Breit-Wigner distribution with a delta function [77], yielding two independent processes: the production and decay of the Higgs boson.

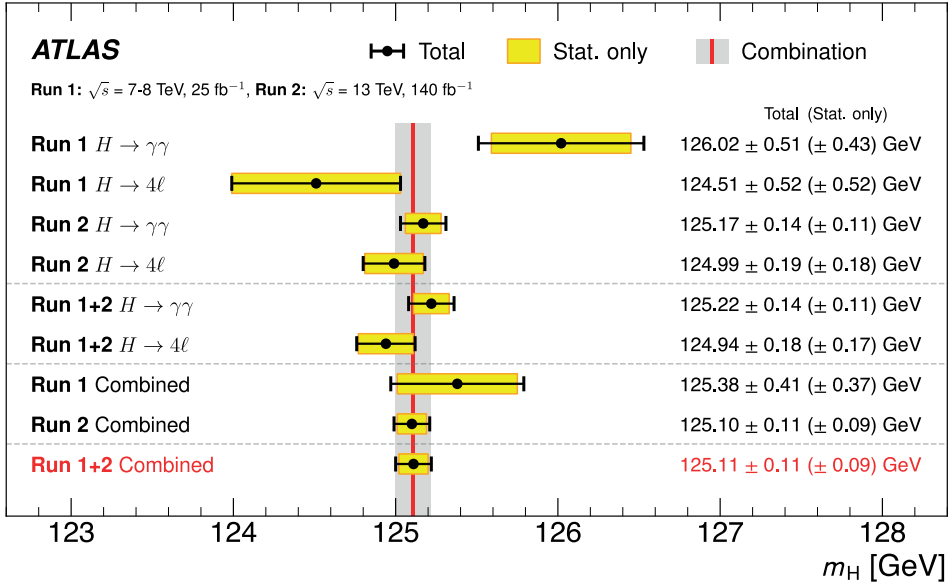
Let us use the process central to this thesis; a Higgs boson produced via the collision of two protons ( $p$ ) and its decay into two leptons ( $\ell$ ), comprised of electrons ( $e$ ), muons ( $\mu$ ), and two neutrinos ( $\nu$ ), via a pair of charged vector bosons ( $W^\pm$ ). This process is typically denoted as  $pp \rightarrow H \rightarrow \ell\nu\ell\nu$  and shorthand for  $pp \rightarrow H \rightarrow W^-W^+ \rightarrow \ell^+\nu\ell^-\bar{\nu}$ .

The total cross-section is written as

$$\begin{aligned}\sigma(pp \rightarrow H \rightarrow \ell\nu\ell\nu) &= \frac{g_{pp \rightarrow H}^2 g_{H \rightarrow \ell\nu\ell\nu}^2}{(s - m_H^2)^2 + m_H^2 \Gamma_H^2} \text{ (NWA } m_H \gg \Gamma_H) \\ &\propto \frac{g_{pp \rightarrow H}^2 g_{H \rightarrow \ell\nu\ell\nu}^2}{m_H \Gamma_H} \delta(s - m_H^2).\end{aligned}$$

This delta function ensures if the Higgs boson is produced on its resonance mass, so called on-shell, it leaves the relation

$$\sigma_{pp \rightarrow H \rightarrow \ell\nu\ell\nu}^{\text{on-shell}} = |M(pp \rightarrow H)|^2 \cdot \frac{\Gamma(H \rightarrow \ell\nu\ell\nu)}{m_H \Gamma_H}.$$



**Figure 3.1:** Summary of Higgs boson mass measurements from individual  $H \rightarrow \gamma\gamma$  and  $H \rightarrow ZZ^* \rightarrow 4\ell$  analyses along with their combinations in Run-1 and Run-2 [75].

In this relation,  $M(pp \rightarrow H)$  is the matrix element describing a Higgs boson produced on-shell,  $\Gamma(H \rightarrow \ell\nu\ell\nu)$  is the partial width Higgs decay to the final state  $\ell\nu\ell\nu$  and  $\Gamma_H$  is total width of the Higgs boson. Using this equation the branching ratio to a final state  $\ell\nu\ell\nu$  is defined as  $BR(H \rightarrow \ell\nu\ell\nu) = \Gamma(H \rightarrow \ell\nu\ell\nu)/\Gamma_H$ .

According to Fermi's golden rule [78], the decay width of a Breit-Wigner resonance

$$\Gamma = \frac{2\pi}{\hbar} |M_{H \rightarrow \ell\nu\ell\nu}|^2 \frac{dN}{dE_f},$$

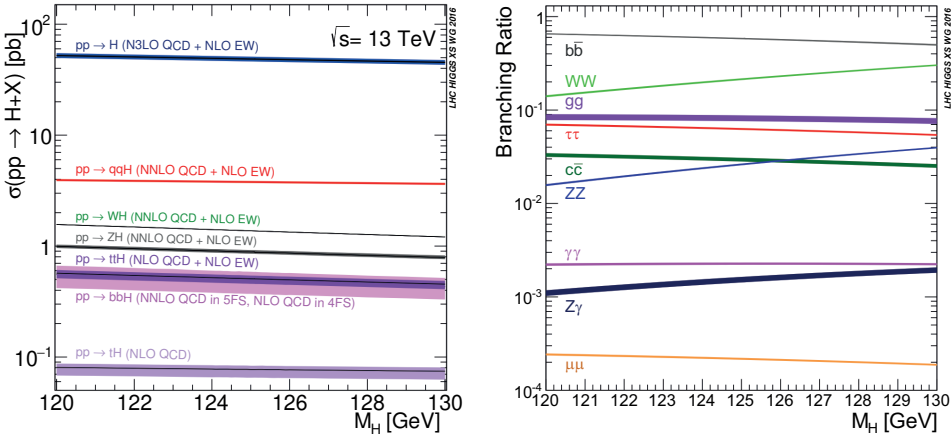
relates the width  $\Gamma$  directly to the matrix element, the phase-space of the final state, the density of the final states  $dN/dE_f$  and identifies  $\Gamma$  as the rate of decay [79]. The spontaneous emission decay law, without the presences of external fields, is given by an exponential decay

$$\frac{dN}{dt} = -\Gamma N(t).$$

After integration, this results in  $N(t) = N_0 e^{-\Gamma t}$  and contains the average lifetime  $\tau = 1/\Gamma$ . The parameter  $\tau$  is the average time it takes for a resonance to decay. Via the use of this relation and  $\Gamma = 4.1\text{ MeV}$ , the SM Higgs boson average lifetime is predicted to be  $\tau_H [\text{GeV}^{-1}] = 1/\Gamma_H [\text{GeV}^{-1}] = \hbar/\Gamma_H [Js/J] = 1.6 \times 10^{-22}\text{ s}$ . A measurement of the Higgs boson width can therefore be interpreted as a constraint on its lifetime and used to probe potential couplings to new particles.

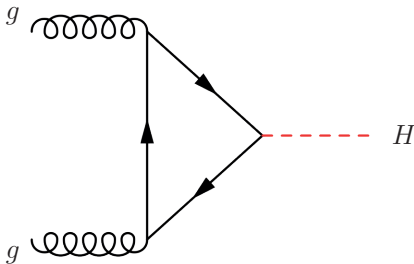
### 3.2 Higgs production modes and decay channels

In Fig. 3.2, the production cross-sections and branching ratios of the SM Higgs boson used for the lifetime calculation are shown as function of the mass of the Higgs boson. The production cross-sections are calculated at  $\sqrt{s} = 13$  TeV, the Run-2 centre-of-mass energy of the LHC.

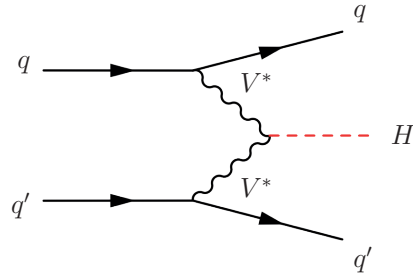


**Figure 3.2:** The dominant Higgs production mode cross-sections (left) [80] and branching ratios (right) [76] as function of the Higgs mass.

The Feynman diagrams of the main Higgs boson production modes are denoted as ggF, VBF,  $VH$ ,  $t\bar{t}H$  and  $tH$  and are shown in Figs. 3.3 and 3.4. The dominant production mode for the Higgs, with  $m_H = 125$  GeV, is the ggF or gluon-gluon fusion process. In this process, two gluons originating from the colliding protons form a quark loop which then annihilate to form the Higgs boson. The main quark loop contributing to this process is the top quark loop, because it has the largest mass.



The ggF production mode.



The VBF production mode.

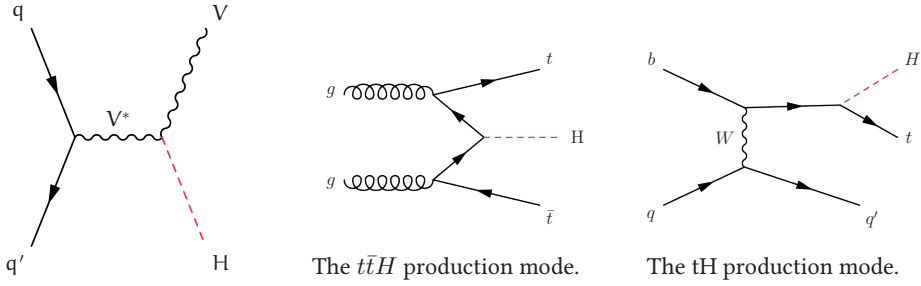
**Figure 3.3:** The Feynman diagrams illustrating the production modes of the Higgs boson analysed and measured in this thesis.

The second largest production mode is the VBF or the vector boson fusion process: in



this process two quarks from the colliding protons radiate  $W$  or  $Z$  bosons, which then fuse to form a Higgs boson in association with two quark-initiated jets [76, 80].

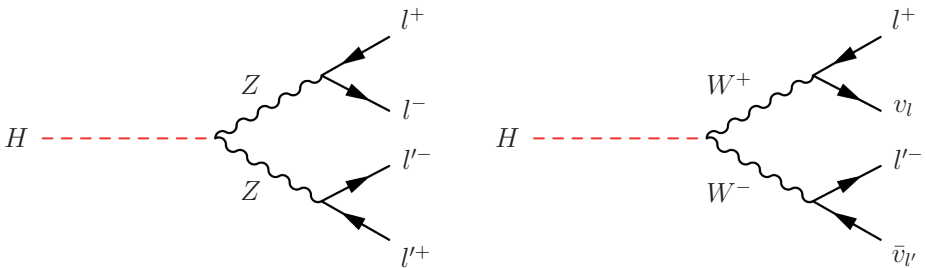
The third largest production mode is the  $VH$  or vector boson associated production [76, 80]: in this process, an s-channel  $W$  or  $Z$  boson is produced which radiates a Higgs boson. Last are the  $t\bar{t}H$  and  $tH$  processes, which produce one or two top quarks in association with a Higgs boson via the top-Yukawa ( $Y_t$ ) coupling.



The  $VH$  production mode.

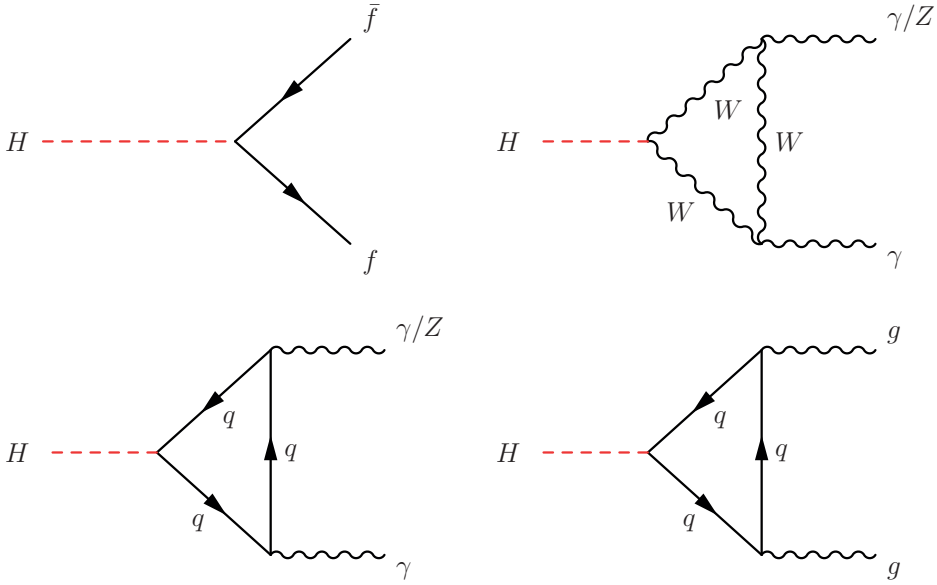
**Figure 3.4:** The Feynman diagrams illustrating the less dominant  $VH$ ,  $t\bar{t}H$  and  $tH$  production modes of the Higgs boson.

A Higgs boson can decay into a variety of final states because it couples to all massive particles in the SM [76, 80]. The top quark, the most massive known SM particle, has the largest coupling to the Higgs boson. However, since a top quark pair is much heavier compared to the on-shell Higgs mass, this decay is suppressed. Therefore, the most dominant decay is the  $H \rightarrow b\bar{b}$  decay mode, since the bottom quark is the second heaviest quark in the SM and the total mass of two  $b$ -quarks, unlike two top quarks, is less than 125 GeV. After these decay modes, the  $H \rightarrow WW^*$  and  $H \rightarrow ZZ^*$  decays depicted in Fig. 3.5 follow, in which one of the  $W$  or  $Z$  bosons can decay off-shell (indicated by the asterisk), making these decay modes kinematically suppressed.



**Figure 3.5:** The Feynman diagrams illustrating the  $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$  and  $H \rightarrow ZZ^* \rightarrow 4\ell$  decay modes of the Higgs boson, in which one of the  $W$  or  $Z$  bosons can decay off-shell

The other kinematically allowed modes, which decay into fermion and boson pairs, are shown in Fig. 3.6. The Higgs can decay into all massive leptons and, via loops through intermediate quarks or  $W$  bosons, into  $\gamma\gamma$ ,  $gg$ , or  $Z\gamma^*$ ; these are known as loop-induced decays.



**Figure 3.6:** The Feynman diagrams illustrating the  $H \rightarrow f\bar{f}$  decay mode, which includes all massive leptons and quarks. The  $H \rightarrow \gamma\gamma/Z\gamma$  decay mode via an intermediate  $W$ -loop and the  $H \rightarrow \gamma\gamma/Z\gamma$  or the  $H \rightarrow gg$  decay mode via an intermediate quark loop.

### 3.3 Measurements of Higgs boson cross-sections

Since the discovery of the Higgs boson, various techniques have been developed to measure its properties. As the LHC Run-2 dataset grew to the  $140 \text{ fb}^{-1}$  used in this thesis, the total number of produced Higgs bosons increased, enhancing the dataset's statistical strength. This improvement allows for more precise measurements, simultaneous measurement of multiple parameters, and measurements of previously statistically limited parameters. The measurements in this thesis are categorized into three main types: signal strength measurements, fiducial cross-section measurements, and Simplified Template Cross-section (STXS) measurements.

#### 3.3.1 Signal strength measurements

Statistically limited analyses, for example rare Higgs process measurements and the measurement discussed in Chapter 10, primarily focus on discoveries of processes or assessing a production rate. These measurements concentrate their precision on a single parameter, the inclusive signal strength.

The signal strength ( $\mu$ ) for Higgs boson production and decay in the process  $pp \rightarrow H \rightarrow \ell\nu\ell\nu$ , is defined as

$$\mu_{pp \rightarrow H \rightarrow \ell\nu\ell\nu} = \frac{(\sigma_{pp \rightarrow H} \cdot B_{H \rightarrow \ell\nu\ell\nu})^{\text{observed}}}{(\sigma_{pp \rightarrow H} \cdot B_{H \rightarrow \ell\nu\ell\nu})^{\text{SM}}}.$$

The signal strength compares the measured signal yield to the yield predicted by the theory in the analysis regions and corrects for the efficiency induced by the detector or the measurement. Since only one parameter is extracted from the fit, this method provides limited information about any BSM effects that might occur in less populated regions of the measured phase-space.

### 3.3.2 Fiducial differential cross-section measurements

As BSM physics is not limited to changing the production rate of SM processes, but can also alter the kinematic properties of the final state particles, a measurement technique that provides more information about the kinematic properties is used in [Chapter 13](#).

The fiducial differential cross-section measurements are a significantly more detailed version of the signal strength measurements and provide information about the total rate and the kinematics of the process of interest. The fiducial phase-space, is a phase-space at particle-level designed to correspond closely to what is experimentally accessible by the detector and reconstruction/identification algorithms, which allows the measurement to be as model-independent as possible. The result of this approach is a differential measurement of the cross-section as a function of, for example, the invariant mass of the final state particles. This technique is typically more flexible for reinterpreting the results in the context of BSM physics because of looser selections and can provide more information about possible deviations in the tail of distributions.

The fiducial differential cross-section result can be transformed into a differential cross-section result by correcting for detector acceptance and efficiency; this extrapolates the result from the reconstructed analysis phase-space into the phase-space at the particle-level. This procedure, called unfolding [\[81, 82\]](#), is useful as it allows for a direct comparison between the measured differential cross-section and arbitrary theoretical predictions or BSM models in the future without the need of additional detector and reconstruction simulation.

### 3.3.3 The Simplified Template Cross-Section framework

In order to use combination of the strengths of the aforementioned analysis technique and efficiently probe potential kinematic deviations of the Higgs boson production, a mixture of the signal strength and fiducial differential cross-section approaches is used in the form of the STXS framework [\[76, 83\]](#). This framework combines the strengths of the previously mentioned techniques and is defined through a partition of the particle-level signal phase-space into predefined STXS regions motivated by BSM theory and SM Higgs production. Each of these STXS regions is then assigned an dedicated signal strength at

particle-level. This measurement technique is used in [Chapter 12](#) and combined with the fiducial differential cross-section measurements in [Chapter 13](#).

Since every STXS region is designed to have sensitivity to a different production mode or kinematic property of the process of interest, each orthogonal region essentially functions as a separate inclusive signal strength or cross-section measurement. This enables a more model-independent representation of the measurement results because of the added differential information.

The framework also retains the possibility for each analysis to use their own measurement techniques and reconstructed objects most optimal for their channel, while supplying the ability to easily combine results with other Higgs analysis on particle-level in the form of the STXS region signal strengths. In contrast to standard signal strength measurements, where combination is only possible on the reconstruction level because of the lack of predefined and agreed upon particle-level signal categories. Another useful feature of the STXS framework is the ability to evolve with the amount of data or analysis sensitivity used in the measurement, allowing for finer splits in the kinematic observables if the sensitivity is deemed sufficient. Similarly, when the sensitivity is not sufficient, the regions can be merged to increase the statistical precision.

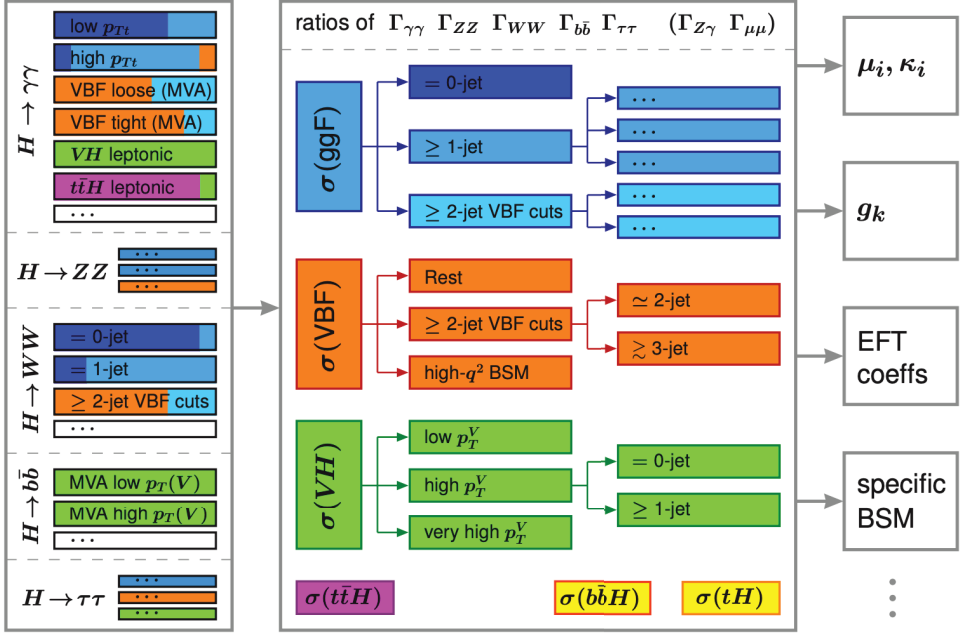
The STXS regions are orthogonal to each other in order to be used in a simultaneous profile likelihood fit during a combination of measurements. Specifically, they are defined using particle-level information of the Higgs boson and any of its associated particles. In this process, the particle-level phase-space is first split into Higgs production modes:  $ggF$ ,  $q\bar{q}H$  (VBF),  $VH$ ,  $t\bar{t}H$ , and  $tH$ . Following the production mode categorisation, further splits are performed based on the kinematics of the produced particles. The regions are chosen to maximize sensitivity to potential deviations from the SM while avoiding large theoretical uncertainties. Additionally, all production topologies are required to have a Higgs boson rapidity of  $|y_H| < 2.5$ , which corresponds to the experimental acceptance of the ATLAS detector.

The number of signal events in each STXS category  $k$  can be expressed as

$$n_k^{\text{signal}} = \mathcal{L}_k \sum_{i,f,k'} (\sigma \times B)^{i,f,k'} (A \times \epsilon)_k^{i,f,k'},$$

where the sum over  $i$  runs over the several production modes included by the input analysis. The superscript  $f$  denotes the final states ( $WW$ ,  $ZZ$ ,  $\gamma\gamma$ ,  $\tau\tau$ ,  $b\bar{b}$ ) and  $k'$  is the index that represents a particle-level STXS category. The parameter  $\mathcal{L}_k$  is the integrated luminosity of the dataset in reconstructed category  $k$ , and  $(A \times \epsilon)_k^{i,f,k'}$  is the acceptance times efficiency factor in category  $k$ . The product  $(\sigma \times B)^{i,f,k'}$  is the cross-section times branching fraction for production mode  $i$ , final state  $f$  and particle-level STXS region  $k'$ . The number of reconstructed signal events,  $n_k^{\text{signal}}$ , can then be interpreted in terms of either signal strengths, coupling modifiers ( $g_k$ ), (SM-)EFT coefficients ( $c_i$ ), or any other BSM model. When interpretations of STXS measurements are performed with BSM

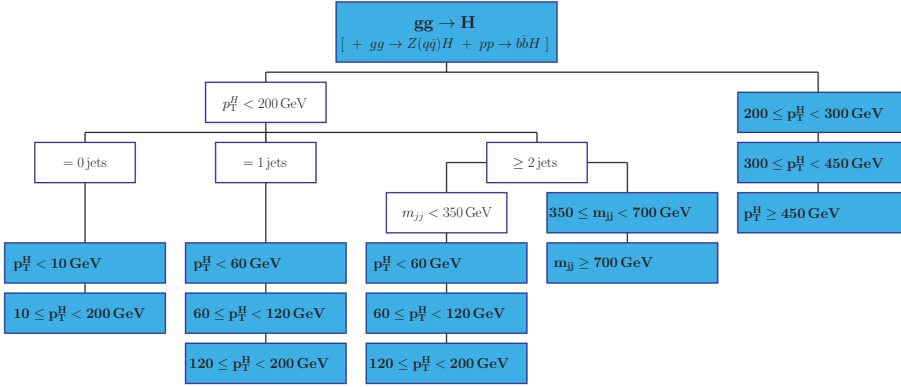
models that significantly alter kinematic distributions, additional acceptance correction factors are calculated to account for changes in the acceptance and signal shape as a function of the model parameters. The STXS methodology is illustrated in Fig. 3.7.



**Figure 3.7:** A schematic overview of the STXS framework and its different sections. On the left: the analyses are illustrated. In the middle: the particle-level STXS regions are shown. To the right: the various possible theory interpretations are indicated. The particle-level phase-space is split into orthogonal STXS regions which are based on different production modes and kinematic properties of the Higgs boson. The signal events in each region are then used to either measure the signal strength, coupling modifiers, EFT coefficients or BSM models while allowing each analysis to use their own measurement techniques on reconstructed objects [76].

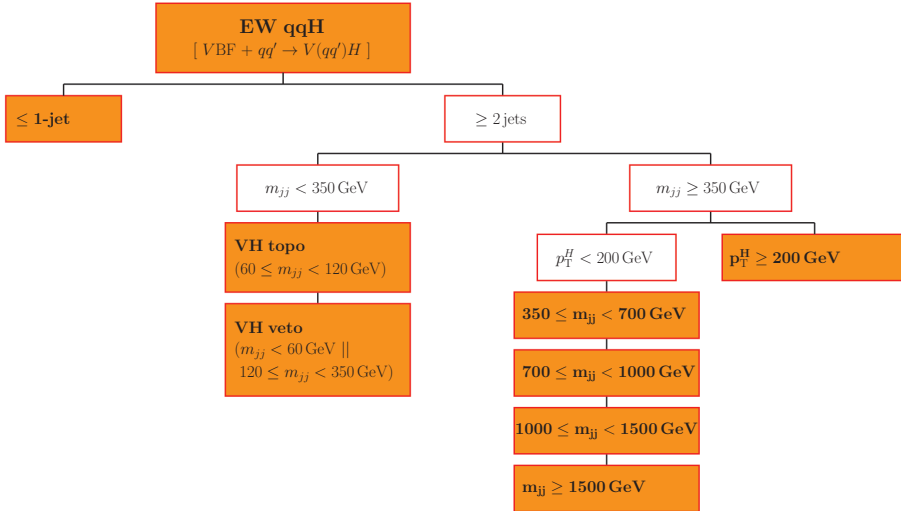
The STXS categories used in this thesis are defined as follows:

- $gg \rightarrow H$  production mode: This category also includes  $gg \rightarrow ZH$  and  $gg \rightarrow b\bar{b}H$  production because of the gluon initiation. The category is first separated into  $p_T^H < 200$  GeV and  $p_T^H > 200$  GeV regions. The high- $p_T^H$  regions are sensitive to any BSM effects which the Higgs boson may recoil against; the low- $p_T^H$  regions are split into jet bins ranging from 0 to 2 or more jets and further split into several  $p_T^H$  and  $m_{jj}$  regions to provide a more differential measurement of the Higgs boson production. This split is illustrated in Fig. 3.8.
- EW  $pp \rightarrow q\bar{q}H$  production mode: Because, at higher-order in QED, the VBF and  $VH$  with hadronically decaying  $V$  bosons are produced through the same process, they are included in the same STXS bin named EW  $q\bar{q}H$ . The process is split into a 2-jet and less-than-2-jet region, to account for the cases where jet is outside of the detector acceptance or mis-identified. The  $q\bar{q}H$  process is further split at



**Figure 3.8:** Definition of the STXS measurement regions for the ggF production mode. The ggF production mode is split into different regions via intermediate into the final regions [6].

the analysis level into VBF and  $VH$  regions using  $m_{jj} < 350$  GeV. The low- $m_{jj}$  region is enriched with  $VH$ -like topologies; high- $m_{jj}$  regions are enriched with VBF. A further split of  $p_T^H > 200$  GeV is employed in the VBF high- $m_{jj}$  region to enhance sensitivity to BSM effects. The EW  $q\bar{q}H$  split is illustrated in Fig. 3.9.

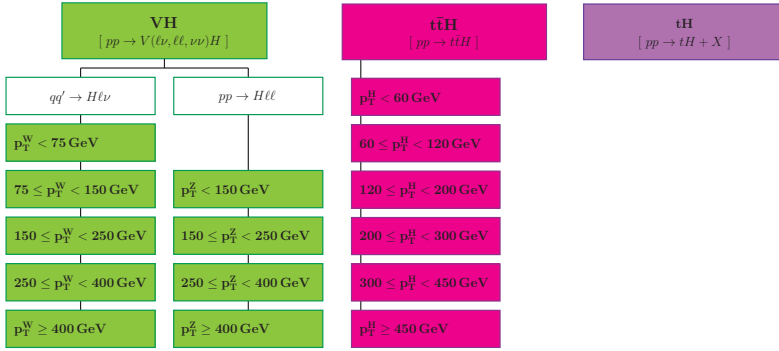


**Figure 3.9:** Definition of the STXS measurement regions for the  $q\bar{q}H$  production mode. The  $q\bar{q}H$  production mode, indicated by colour, is split into different regions via intermediate boxes into the final lower boxes [6].

- $pp \rightarrow V(\ell\nu, \ell\ell, \nu\nu)H$  production mode: The  $VH$  production with leptonically decaying V bosons has its own STXS regions because it is more easily reconstructed and sensitive to BSM effects. The regions are further split into  $WH$  and  $ZH$  production, in order to measure these processes separately. Instead of splitting the regions based on  $p_T^H$ , the regions are now split in  $p_T^V$ , because this gives a less

biased proxy of the Higgs boson boost compared to  $p_T^H$ . Additionally, it allows the usage of the more numerous  $H \rightarrow b\bar{b}$  decays, without relying on its reconstruction for the discriminant variable.

- $t\bar{t}H$  and  $tH$  production modes: The Higgs boson production in association with a  $t\bar{t}$  pair ( $t\bar{t}H$ ) is split into kinematic bins of the transverse momentum of the Higgs boson ( $p_T^H$ ). This differential information is sensitive to any BSM effects, while simultaneously not relying on the reconstruction of the top quarks. The  $tH$  production mode is evaluated inclusively, as the amount of data does not allow for a differential measurement yet. The  $VH$ ,  $t\bar{t}H$ , and  $tH$  categorisation is illustrated in Fig. 3.10.



**Figure 3.10:** Definition of the STXS measurement regions for the  $VH$ ,  $t\bar{t}H$ , and  $tH$  production modes. Each production mode, indicated by colour, is split into different regions via intermediate boxes into the final lower boxes [6].

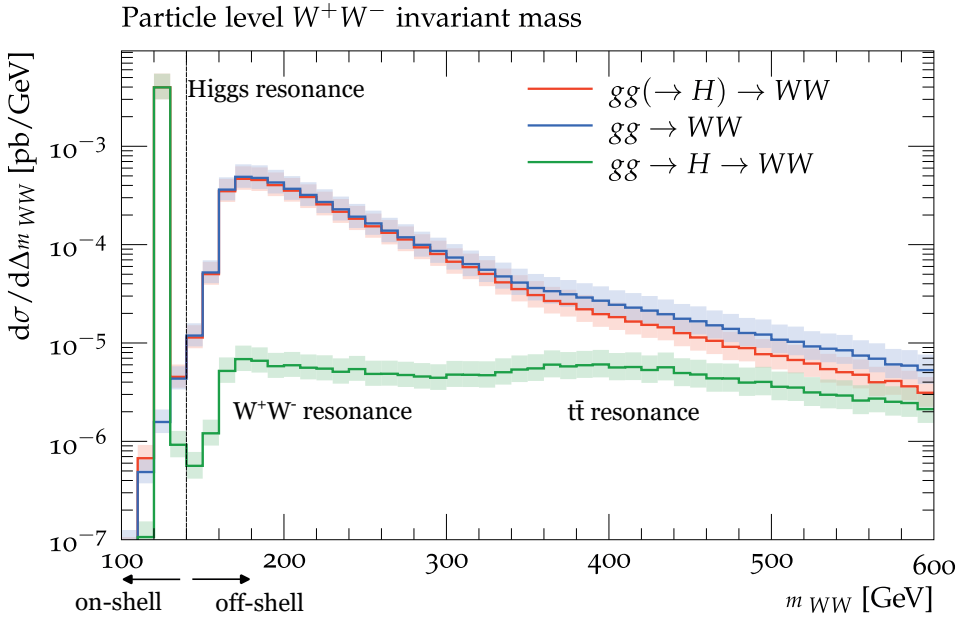
### 3.4 Off-shell Higgs physics

Since any deviation from the SM Higgs width prediction hints towards physics beyond the SM, measuring it has been an object of study ever since its discovery. Unfortunately, the resolution of the ATLAS detector is almost three orders of magnitude worse than  $\Gamma_H$ , making it impossible to measure the width directly from the mass distribution of its decay products with competitive precision. Alternative approaches [84–86] have been suggested and applied in recent years to improve the sensitivity towards this Higgs boson property. One of these approaches uses the off-shell Higgs production and is applied in Chapter 10.

#### 3.4.1 Off-shell couplings and interference

The invariant mass spectrum of Higgs boson production peaks at 125 GeV, where the Higgs is produced on-shell, and decreases significantly in the off-shell region, defined as  $m_H > 140 \text{ GeV}$ , the further it is produced from its resonance mass, as shown in Fig. 3.11. However, the Higgs production rate increases again when its decay products the vector bosons become on-shell, forming a second peak at  $m_{ZZ} \sim 182 \text{ GeV}$  or  $m_{WW} \sim 161 \text{ GeV}$ ,

depending on the final state of interest. Furthermore, as the dominant ggF production mode involves mostly top quarks in its loop, the probability increases once more when both top quarks become on-shell, satisfying  $m_{t\bar{t}} \sim 346$  GeV, which creates a third peak. The  $WW/ZZ$  and  $t\bar{t}$  resonances therefore extend the Higgs production cross-section well into the high-mass regime, resulting in a significant amount of off-shell Higgs bosons. Due to the threshold effects at the  $m_{VV}$  and  $m_{t\bar{t}}$  resonances, the off-shell region contributes  $\mathcal{O}(15\%)$  to the total Higgs boson production cross-section. The behaviour of the Higgs boson production cross-section as a function of the invariant mass of the  $WW$  decay products.



**Figure 3.11:** The ggF  $H \rightarrow WW$  cross-section decaying into  $e\nu\mu\nu$ , as calculated by Sherpa at LO, as a function of the invariant mass of the  $WW$  decay products. The cross-section of the  $gg(\rightarrow H) \rightarrow WW$  process is shown in the on-shell and off-shell region in red. The same is shown for the  $gg \rightarrow WW$  and  $gg \rightarrow H \rightarrow WW$  in blue and green, respectively. The off-shell region is defined as  $m_{WW} > 140$  GeV [6] and the on-shell as  $m_{WW} < 140$  GeV. The shaded bands represent the scale and PDF uncertainties added in quadrature.

The cross-section dependence on  $\Gamma_H$  becomes negligible when the Higgs is sufficiently off-shell [87]. Therefore, when making the off-shell assumption, i.e.  $(m_{WW} - m_H) \gg \Gamma_H$ , the mass difference between the decay products and the on-shell mass is much greater compared to the width of the Higgs boson resonance Breit-Wigner peak. This leads to a cross-section that is independent of the Higgs boson width, for an off-shell produced Higgs boson decaying into  $WW$  this is:

$$\sigma(pp \rightarrow H^* \rightarrow WW) \propto \frac{g_{pp \rightarrow H}^2 g_{H \rightarrow WW}^2}{(s - m_H^2)^2}.$$

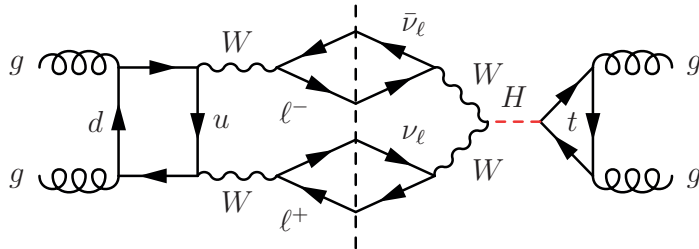


In the on-shell region, the cross-section scales by  $\Gamma_H$  in the denominator. Therefore, by measuring both the on-shell and off-shell Higgs boson production, and evaluating their ratio, a handle is provided for constraining the Higgs boson width and its average lifetime.

### 3.4.2 The diboson continuum

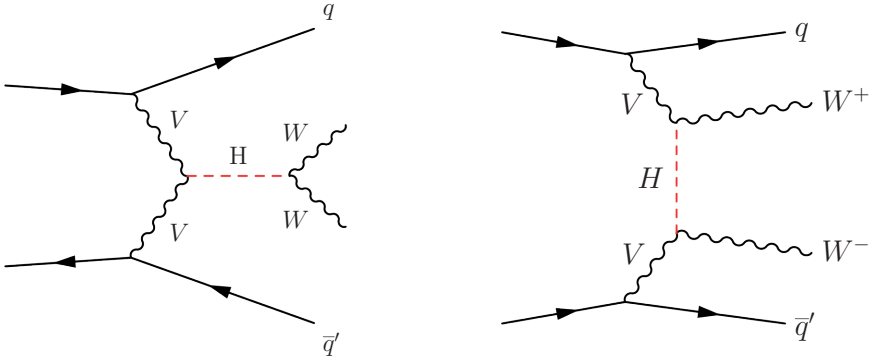
The diboson continuum,  $pp \rightarrow W^+W^- + X$ , plays an important role in the off-shell region. Since the process has the same initial and final state as the mechanism producing Higgs bosons that decay via vector bosons, the two processes interfere. This interference can be constructive or destructive depending on the phase of the amplitude; in the case of the SM Higgs boson, the interference is destructive. This is caused by the Higgs boson's positive parity and the electroweak diboson continuum's negative parity; it has been experimentally verified that weak interactions do not conserve parity [88], causing the interference to be destructive due to the sign of the  $W$  boson propagator along with its VA nature [30].

Therefore, the total cross-section of  $gg \rightarrow (H^*) \rightarrow WW$ , in which two  $W$  bosons are produced either with or without the interplay of a Higgs boson, is reduced [76] due to the destructive interference. The topology in which this thesis explores the off-shell nature of the Higgs boson is the  $W^+W^-$  channel. The leading contributing Feynman diagrams involving Higgs bosons decaying into  $W^+W^-$ , and the interfering background, are shown in Fig. 3.12 for the  $ggF$  production mode.

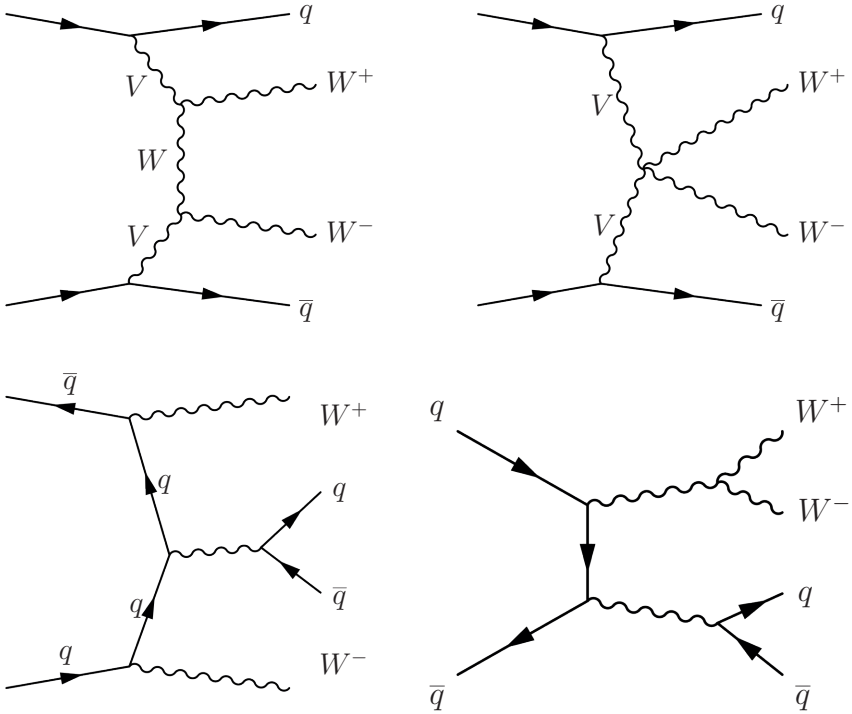


**Figure 3.12:** The LO Feynman diagram involving a  $W$  boson pair produced by  $gg$ -induced closed fermion loop (left) and the mirrored Higgs signal which also produces a  $W$  boson pair (right) and features an identical initial and final state. The dotted black line mirrored the right Feynman diagram in order to illustrate interference occurring between these two diagrams in the  $gg$ -channel.

The VBF (or EW) off-shell Higgs production process, depicted in Fig. 3.13, has a Higgs resonance with a less distinct  $W^+W^-$  peak because of the energy carried away by the two radiating quarks in the process. This process interferes destructively with vector boson scattering and quark annihilation processes shown in Fig. 3.14.



**Figure 3.13:** The LO Feynman diagrams of Higgs bosons production by electroweak vector boson interactions.



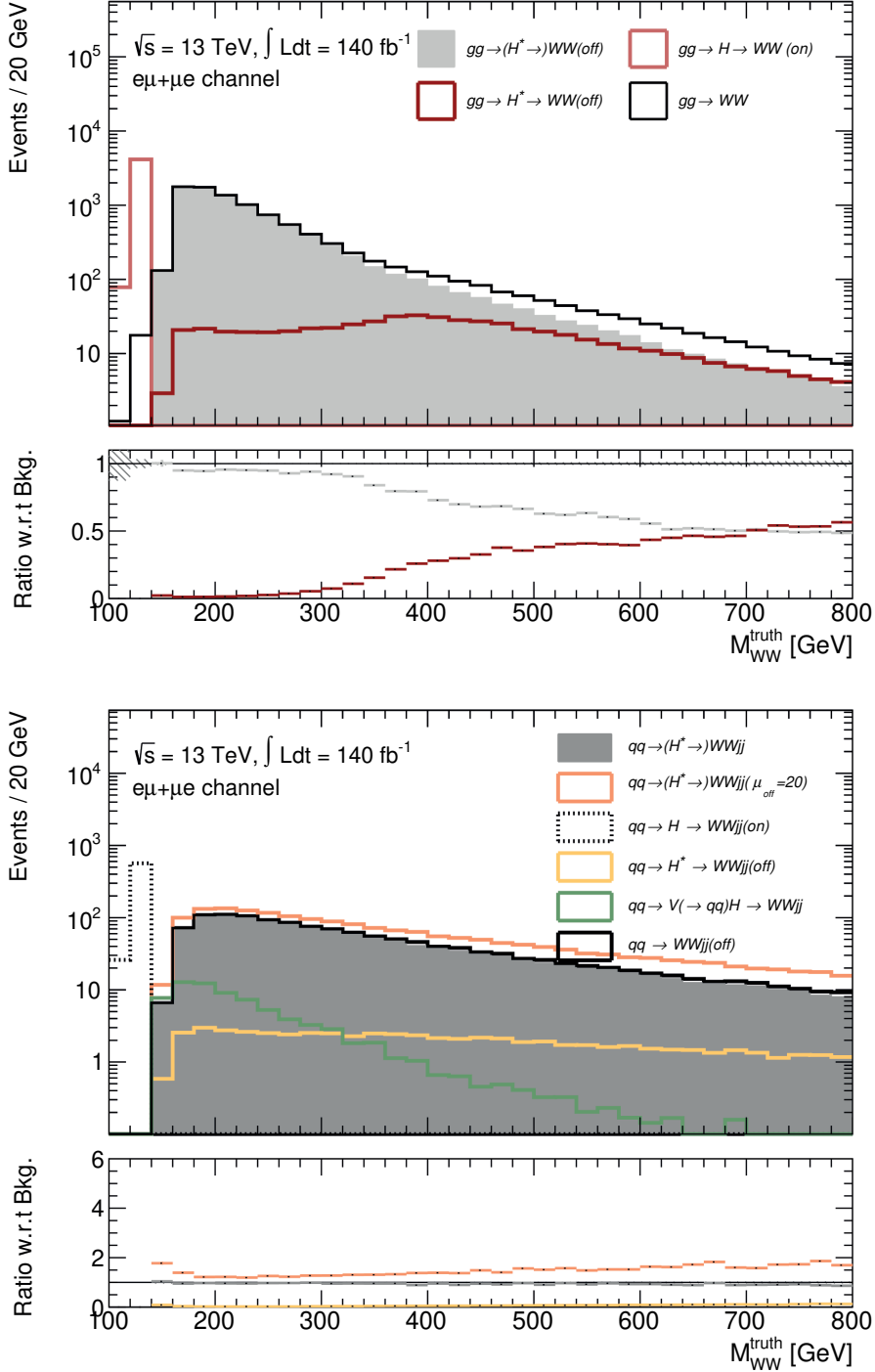
**Figure 3.14:** The LO Feynman diagrams of vector boson scattering and quark annihilation background processes that interfere with electroweak Higgs boson production decaying as  $H \rightarrow WW$ .

The  $m_{WW}^{\text{truth}}$  system kinematics are reconstructed using truth (particle-level) information (before showering and hadronization) and stored for later use during the analysis. The Higgs boson contribution and its interfering background are split into their off-shell and on-shell components using the  $m_{WW}^{\text{truth}}$  system information during the analysis. The off-shell and on-shell components, used in the analysis, are defined as  $m_{WW}^{\text{truth}} \geq 140$  GeV and  $m_{WW}^{\text{truth}} < 140$  GeV, respectively. The  $m_{WW}^{\text{truth}}$  selection of 140 GeV is chosen in order to ensure the off-shell Higgs bosons are sufficiently high-mass such that the  $\Gamma_H$  parameter has no influence on its cross-section.

To increase the efficiency of event generation, an event-level filter is applied to the VBF and ggF samples. For the ggF samples, an  $m_{\ell\ell} > 45$  GeV selection is used, while for the VBF samples, a  $p_T^{\text{lead. lep.}} > 10$  GeV selection is required. Tab. 3.1 contains the inclusive cross-sections of the used samples. The ggF and VBF production mode line shapes of the  $m_{WW}$  invariant mass, used in Chapter 10, along with their interfering background, are shown in Fig. 3.15.

**Table 3.1:** Calculated cross-sections of the  $qq \rightarrow WW + jj \rightarrow \ell\nu\ell\nu + jj$  and  $gg \rightarrow (H^*) \rightarrow WW \rightarrow \ell\nu\ell\nu$  different-flavour processes.

| Process                                                                                        | Cross-section (pb)               |                                   |
|------------------------------------------------------------------------------------------------|----------------------------------|-----------------------------------|
|                                                                                                | Order: up to LO ( $\alpha_s^2$ ) | Order: up to NLO ( $\alpha_s^3$ ) |
| $gg \rightarrow (H^*) \rightarrow WW \rightarrow \ell\nu\ell\nu$                               | 79.8                             | 158.4                             |
| $gg \rightarrow WW \rightarrow \ell\nu\ell\nu$                                                 | 44.1                             | 75.6                              |
| $gg \rightarrow H^* \rightarrow WW \rightarrow \ell\nu\ell\nu$                                 | 40.1                             | 90.3                              |
| Order: up to LO( $\alpha_{\text{EW}}^6$ )                                                      |                                  |                                   |
| $qq \rightarrow (H \rightarrow)WW \rightarrow \ell\nu\ell\nu jj (\mu_{\text{off-shell}} = 1)$  | 0.176                            |                                   |
| $qq \rightarrow (H \rightarrow)WW \rightarrow \ell\nu\ell\nu jj (\mu_{\text{off-shell}} = 20)$ | 0.217                            |                                   |
| $qq \rightarrow WW \rightarrow \ell\nu\ell\nu jj$                                              | 0.101                            |                                   |



**Figure 3.15:** The ggF (top) and VBF (bottom) production mode along with its interfering background. The Higgs boson contribution and its interfering background are split into their off-shell ( $m_{WW} \geq 140$  GeV) and on-shell ( $m_{WW} < 140$  GeV) components. Along with the  $m_{WW}$  distribution in the top pad, the ratio with respect to the background-only sample is calculated in the lower pad.

### 3.4.3 The lifetime measurement

In the years 2013 and 2014, first for the  $ZZ$  and later for the  $WW$  channel, it was proposed to combine the signal strength measurement of the off-shell analysis with a signal strength measurement of the on-shell Higgs production in order to constrain the total Higgs width ( $\Gamma_H$ ) [84, 85].

The rate of  $WW$  production mediated by an on-shell Higgs  $H$  boson depends on the Higgs boson's couplings to particles in its production and decay,  $g_{\text{prod}}(m_H)$  and  $g_{\text{decay}}(m_H)$ , as well as its total width,  $\Gamma_H$ . In contrast, the rate of  $WW$  production mediated by an off-shell Higgs boson  $H^*$  is independent of the total width but still depends on the couplings  $g_{\text{prod}}(\hat{s})$  and  $g_{\text{decay}}(\hat{s})$ , where  $\hat{s}$  represents the virtuality of the Higgs boson.

Deviations from the SM predictions for  $gg \rightarrow H \rightarrow WW$  rates can be expressed in terms of the on-shell and off-shell signal strengths,  $\mu_{\text{on-shell}}$  and  $\mu_{\text{off-shell}}$ , as follows:

$$\begin{aligned}\mu_{\text{on-shell}}^{\text{ggF}} &= (\kappa_{gg \rightarrow H}^{\text{on-shell}})^2 \cdot (\kappa_{H \rightarrow WW}^{\text{on-shell}})^2 \cdot \frac{1}{\kappa_H} \\ \mu_{\text{on-shell}}^{\text{VBF}} &= (\kappa_{H \rightarrow WW}^{\text{on-shell}})^4 \cdot \frac{1}{\kappa_H} \\ \mu_{\text{off-shell}}^{\text{ggF}} &= (\kappa_{gg \rightarrow H}^{\text{off-shell}})^2 \cdot (\kappa_{H \rightarrow WW}^{\text{off-shell}})^2 \\ \mu_{\text{off-shell}}^{\text{VBF}} &= (\kappa_{H \rightarrow WW}^{\text{off-shell}})^4.\end{aligned}$$

In this set of equations, the parameter  $\kappa_H = \Gamma_H / \Gamma_H^{\text{SM}}$ , and the other  $\kappa$  parameters represent coupling strength modifiers with respect to the SM predictions,  $\kappa_i = g_i / g_i^{\text{SM}}$ , where  $i$  is either the  $gg \rightarrow H$  or  $H \rightarrow WW$  coupling. By defining  $R_i = (\kappa_i^{\text{off-shell}})^2 / (\kappa_i^{\text{on-shell}})^2$  the off-shell/on-shell coupling ratios are obtained:

$$R_{gg} \equiv \frac{(\kappa_{gg \rightarrow H}^{\text{off-shell}})^2}{(\kappa_{gg \rightarrow H}^{\text{on-shell}})^2} \quad R_{VV} \equiv \frac{(\kappa_{H \rightarrow WW}^{\text{off-shell}})^2}{(\kappa_{H \rightarrow WW}^{\text{on-shell}})^2}.$$

These coupling ratios are then used to obtain the following relation between the on-shell and off-shell signal strengths:

$$\begin{aligned}\mu_{\text{off-shell}}^{\text{ggF}} &= R_{gg} \cdot R_{VV} \cdot \mu_{\text{on-shell}}^{\text{ggF}} \cdot \frac{\Gamma_H}{\Gamma_H^{\text{SM}}}, \\ \mu_{\text{off-shell}}^{\text{VBF}} &= R_{VV}^2 \cdot \mu_{\text{on-shell}}^{\text{VBF}} \cdot \frac{\Gamma_H}{\Gamma_H^{\text{SM}}}.\end{aligned}$$

Using this parameterization of the on-shell and off-shell signal strengths, the following two interpretations are considered for an off-shell/on-shell combined measurement.

In the first interpretation, the parameter of interest (POI) is the Higgs boson width normalized to the SM prediction,  $\Gamma_H / \Gamma_H^{\text{SM}}$ . In this interpretation, it is assumed that the coupling strengths for on-shell and off-shell processes are equal, i.e.  $R_{gg} = 1$  and  $R_{VV} = 1$ . This assumption reduces the equations relating the on-shell and off-shell

signal strengths

$$\begin{aligned}\mu_{\text{off-shell}}^{\text{ggF}} &= \mu_{\text{on-shell}}^{\text{ggF}} \cdot \frac{\Gamma_H}{\Gamma_H^{\text{SM}}} \\ \mu_{\text{off-shell}}^{\text{VBF}} &= \mu_{\text{on-shell}}^{\text{VBF}} \cdot \frac{\Gamma_H}{\Gamma_H^{\text{SM}}}.\end{aligned}$$

The ratio  $\Gamma_H/\Gamma_H^{\text{SM}}$  is measured and related to the lifetime of the Higgs boson.

A second interpretation can be made, in which the parameters of interest are the  $R_{gg}$  and  $R_{VV}$  ratios. When these parameters are measured, while  $\Gamma_H/\Gamma_H^{\text{SM}}$  is assumed to be 1, the equations relating the on-shell and off-shell signal strengths reduce to

$$\begin{aligned}\mu_{\text{off-shell}}^{\text{ggF}} &= R_{gg} \cdot R_{VV} \cdot \mu_{\text{on-shell}}^{\text{ggF}} \\ \mu_{\text{off-shell}}^{\text{VBF}} &= R_{VV}^2 \cdot \mu_{\text{on-shell}}^{\text{VBF}}.\end{aligned}$$

By measuring the  $R_{gg}$  and  $R_{VV}$  ratios, the Higgs boson coupling dependence,  $g_{\text{prod}}(\hat{s})$  and  $g_{\text{decay}}(\hat{s})$ , on the virtuality of the Higgs boson can be studied. A ratio deviating from 1 would indicate BSM-like change in the virtuality dependence of the Higgs boson couplings and point towards the presence of BSM physics.

#### 3.4.4 Modelling off-shell Higgs processes and interfering backgrounds

Since the ggF and VBF production modes of the off-shell Higgs process are measured in this thesis using the  $H^* \rightarrow WW \rightarrow \ell\nu\ell\nu$  decay, detailed modelling and thorough understanding of the off-shell signal and its interfering backgrounds is needed. Due to destructive interference,  $H^* \rightarrow WW$  production cannot be studied without considering the background and should be treated as a single  $WW$  production method.

In this thesis, the physical process of producing  $WW$  bosons is referred to as the SBI process, which includes the signal ( $gg \rightarrow H^* \rightarrow WW \rightarrow \ell\nu\ell\nu$ ), the background ( $gg \rightarrow WW \rightarrow \ell\nu\ell\nu$ ), and their interference. When referring to  $WW$  boson production solely through Higgs mediation, it is the Signal (S) process. The following sections detail the modelling of the off-shell Higgs signal (S), the interfering backgrounds (B), and their interference (I).

#### Gluon-gluon fusion Higgs production

The dominant off-shell Higgs production mode, similar to the on-shell mode, is the ggF production. In this mode,  $WW$  bosons are produced through gluon interactions with a closed fermion loop, which then radiates an intermediate Higgs boson or  $WW$  pair directly. The cross-section of the ggF SBI process is defined by adding the signal, background, and interference cross-sections together. This allows the cross-section of the ggF SBI process to be expressed as:

$$\sigma_{\text{SBI}}^{\text{SM}} = \sigma_{gg \rightarrow H^* \rightarrow WW} + \sigma_{gg \rightarrow WW} + \sigma_{\text{interference}}.$$

This signal strength  $\mu_{\text{off-shell}}$  is defined as the off-shell Higgs boson production, decaying into  $WW$ , normalized to the SM prediction. The signal strength scales linearly with the total off-shell Higgs event yield through  $gg \rightarrow H^* \rightarrow WW$  signal process, meaning the interference scales with  $\sqrt{\mu_{\text{off-shell}}}$ . This relationship is evident when expressing the scattering cross-section in terms of the matrix elements of the signal, background, and interference processes:

$$\begin{aligned} \sigma_{\text{SBI}} &\propto |\mathcal{M}_{gg \rightarrow (H^*) \rightarrow WW}|^2 = |\mathcal{M}_{gg \rightarrow H^* \rightarrow WW} + \mathcal{M}_{gg \rightarrow WW}|^2 = \\ &|\mathcal{M}_{gg \rightarrow H^* \rightarrow WW}|^2 + |\mathcal{M}_{gg \rightarrow WW}|^2 \\ &+ 2\text{Re}(\mathcal{M}_{gg \rightarrow H^* \rightarrow WW} \cdot \mathcal{M}_{gg \rightarrow WW}). \end{aligned}$$

When the defined signal strength is inserted into this equation, the following expression is found

$$\begin{aligned} \sigma_{\text{SBI}} &\propto \mu_{gg \rightarrow H^* \rightarrow WW}^{\text{off-shell}} \cdot |\mathcal{M}_{gg \rightarrow H^* \rightarrow WW}|^2 + \mu_{gg \rightarrow WW} \cdot |\mathcal{M}_{gg \rightarrow WW}|^2 \\ &+ 2 \cdot \sqrt{\mu_{gg \rightarrow H^* \rightarrow WW}^{\text{off-shell}}} \cdot \sqrt{\mu_{gg \rightarrow WW}} \cdot \text{Re}(\mathcal{M}_{gg \rightarrow H^* \rightarrow WW} \cdot \mathcal{M}_{gg \rightarrow WW}). \end{aligned}$$

As the goal is to measure the off-shell Higgs signal strength ( $\mu_{\text{off-shell}}$ ), the assumption is made that interfering background is accurately described by the SM. Hence, in order to derive an SBI cross-section as a function of  $\mu_{gg \rightarrow H^* \rightarrow WW}^{\text{off-shell}}$  the assumption  $\mu_{gg \rightarrow WW} = 1$  is made. Now,  $\sigma_{\text{SBI}}(\mu_{\text{off-shell}})$  can be expressed in the form of matrix elements, or in terms of the individual cross-sections:

$$\begin{aligned} \sigma_{\text{SBI}}(\mu_{\text{off-shell}}) &\propto \mu_{gg \rightarrow H^* \rightarrow WW}^{\text{off-shell}} \cdot |\mathcal{M}_{gg \rightarrow H^* \rightarrow WW}|^2 + |\mathcal{M}_{gg \rightarrow WW}|^2 \\ &+ 2 \cdot \sqrt{\mu_{gg \rightarrow H^* \rightarrow WW}^{\text{off-shell}}} \cdot \text{Re}(\mathcal{M}_{gg \rightarrow H^* \rightarrow WW} \cdot \mathcal{M}_{gg \rightarrow WW}). \\ \sigma_{\text{SBI}}(\mu_{\text{off-shell}}) &\propto \mu_{gg \rightarrow H^* \rightarrow WW}^{\text{off-shell}} \cdot \sigma_{gg \rightarrow H^* \rightarrow WW \rightarrow \ell\nu\ell\nu} + \sigma_{gg \rightarrow WW \rightarrow \ell\nu\ell\nu} \\ &+ \sqrt{\mu_{gg \rightarrow H^* \rightarrow WW}^{\text{off-shell}}} \cdot \sigma_{\text{interference}}. \end{aligned}$$

For the off-shell ggF Higgs production process, the cross-sections of the SBI, S and B processes are modelled using Monte Carlo generators and are used to parameterize the SBI process as function of the signal strength.

### Gluon-gluon fusion Higgs simulation

The  $gg \rightarrow (H^*) \rightarrow WW \rightarrow \ell\nu\ell\nu$  SM Higgs signal, its background, and their interference are simulated using the SHERPA 2.2.2 [69] and OpenLoops generator [89–92], with the matrix elements calculated up to one additional parton in the final state. The QCD renormalization and factorization scales are set to  $m_{WW}/2$  and the NNPDF3.0nnlo PDF set [93] is used for the incoming protons.

### 3.4.5 QCD corrections to loop induced off-shell Higgs production

The ggF Higgs process, and in particular the off-shell ggF production mode, is subject to very large higher-order corrections which arise from phase-spaces that are kinematically forbidden at LO but are populated through the emission of hard QCD jets at NLO [94]. These corrections, captured in K-factors, typically increase with the transverse momentum of a single  $W$  boson system. More specifically, this happens at very large transverse momentum  $P_{T,V_1} \approx Q \gg M_W$  with no requirement on the other second vector boson. At LO, because of momentum conservation, the recoil of the first boson is balanced by the second boson, leading to  $p_{T,V_1} = p_{T,V_2}$ . At NLO, the emission of a hard QCD jet in the radiative process  $pp \rightarrow VVj$  can also be balanced by the additional hard jet, with the second vector boson being much softer,  $p_{T,V_2} = \mathcal{O}(M_W)$  [94–96].

Because of these large K-factors, it is important to estimate the higher-order QCD corrections and their uncertainties in the off-shell region for the signal, background, and interference. Unfortunately, as of 2025, obtaining higher-order predictions for these processes is challenging, because NLO corrections for the  $gg \rightarrow H^* \rightarrow WW$  signal with exact quark-mass dependence are not known. This is in contrast with predictions of the on-shell Higgs boson, for which NNLO corrections, assuming the limit  $m_t \rightarrow \infty$ , are available [97–99].

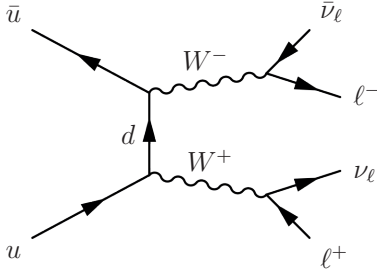
For the off-shell events, the Higgs boson must be treated as virtual without any assumptions on its mass. This requires the calculation of a one-loop amplitude for the LO prediction and a two-loop amplitude for the NLO correction. Additionally, the interference with  $gg \rightarrow WW$  is significant and must be considered in the same calculation. This results in the computation of a process with complex loops and jet activity where the impact of both massive and massless quarks in the box-diagram must be included to arrive at a correct estimation of the yield. This process,  $pp \rightarrow \ell^+ \ell^- \nu_\ell \bar{\nu}_\ell + X$ , contains resonant and non-resonant Feynman diagrams and accounts for off-shell effects and spin correlations [100–103].

#### $qq/qg/gg$ channel and process definitions

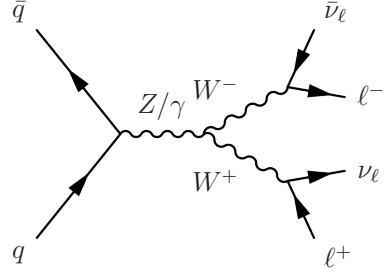
Further attention is spent towards the precise definition of the higher-order corrections used in the off-shell analysis, as the various  $qq/qg/gg$  channels at LO, NLO and NNLO are not easily treated independently and mix in various ways.

Three diagrams of the LO calculation are shown in Figs. 3.16 to 3.18. They are induced by  $q\bar{q}$ -annihilation and include t-channel  $W^+W^-$  topologies, s-channel  $W^+W^-$  topologies, and s-channel Drell-Yan type topologies. Fig. 3.19 shows a LO  $gg$ -diagram induced by a quark loop, where the square of the corresponding amplitude at  $\mathcal{O}(\alpha_s^2)$  is formally also part of the NNLO corrections to the  $q\bar{q}$ -annihilation process.

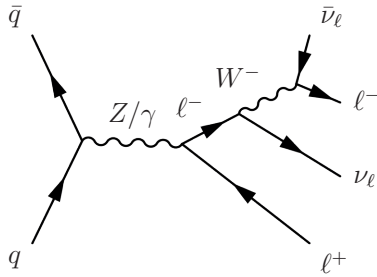




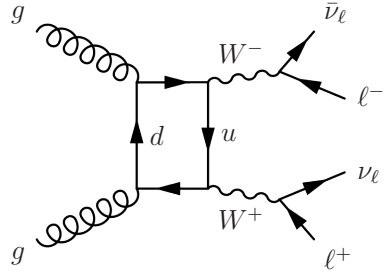
**Figure 3.16:** An example of t-channel  $q\bar{q}$ -annihilation contributing to  $W^+W^-$  production.



**Figure 3.17:** An example of s-channel  $q\bar{q}$ -annihilation contributing to  $W^+W^-$  production.

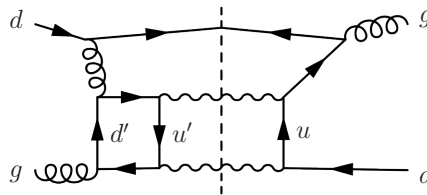


**Figure 3.18:** An example of an Drell-Yan like topology contributing to  $W^+W^-$  production.



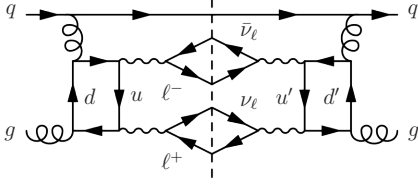
**Figure 3.19:** An example of a  $gg$ -diagram induced by a quark loop contributing to  $W^+W^-$  production.

In addition to  $gg$ -initiated diagrams at  $\mathcal{O}(\alpha_s^3)$ , the NLO corrections also involve  $qg$ -initiated contributions. Not all loop-induced contributions to the  $gg$  channel at  $\mathcal{O}(\alpha_s^3)$  are included, an example is given in Fig. 3.20. This decision is made in order to avoid overlap with quark-induced  $WW$  production; therefore, only contributions where the amplitudes contain a closed fermion loop are considered. This requirement disentangles the mixing quark-induced and gluon-induced interference at NNLO.

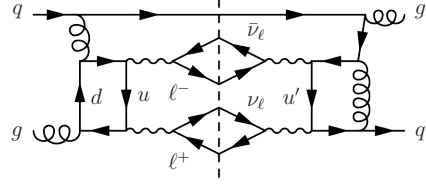


**Figure 3.20:** An example of NNLO diagram interference which is **excluded** in the higher-order calculation. The dotted black line mirrored the right Feynman diagram in order to illustrate interference occurring between these two diagrams in the  $qg$ -channel due to their identical initial and final state but a different internal process ( $u'$  and  $d'$  versus only  $u$ )

A closed loop example included in the corrections is shown in [Fig. 3.21](#), with non-closed loop example, excluded from the correction, shown in [Fig. 3.22](#).

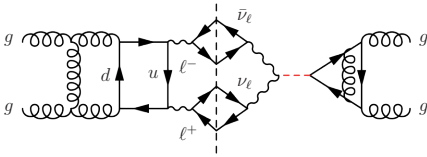


**Figure 3.21:** An example of NNLO contributions **included** in the higher-order calculation. The dotted black line mirrored the right Feynman diagram in order to illustrate interference occurring between these two diagrams in the  $qg$ -channel due to their identical initial and final state but a different internal process ( $u$  and  $d$  versus  $u'$  and  $d'$ ).

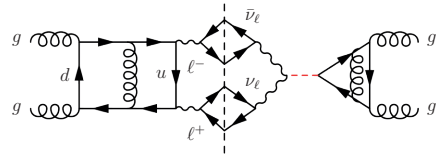


**Figure 3.22:** An example of NNLO diagram interference which is **excluded** in the higher-order calculation. The dotted black line mirrored the right Feynman diagram in order to illustrate interference occurring between these two diagrams in the  $qg$ -channel due to their identical initial and final state but a different internal process ( $u$  and  $d$  versus  $u'$  and  $g$ ).

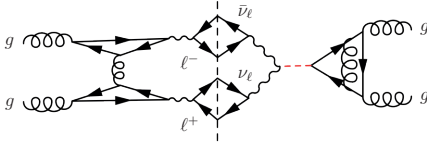
Currently, it is not possible to consistently account for all contributions in a complete  $N^3$ LO calculation. For sufficiently inclusive observables, the impact of missing  $N^3$ LO contributions is expected to fall within the perturbative uncertainties estimated from NNLO scale variations. In contrast, NLO corrections to the loop-induced  $gg$  channel exceed the uncertainties obtained through LO scale variations because of the new  $qg$ -channels entering at this order in perturbation theory [102]. The corrections calculated in this section therefore include all contributions up to  $\mathcal{O}(\alpha_s^2)$  and complete NLO corrections to the loop-induced gluon fusion channel at  $\mathcal{O}(\alpha_s^3)$ . For the  $gg$ -channel, NLO contributions include signal-background interference, illustrated in [Figs. 3.23 to 3.25](#), and in [Fig. 3.26](#) for the  $qg$ -channel.



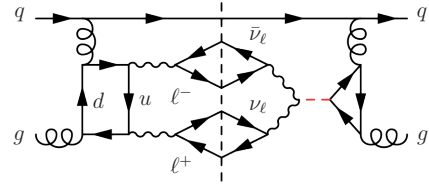
**Figure 3.23:** An example of a two-loop NLO contribution to the  $gg$ -channel with a closed fermion loop. The dotted black line mirrored the right Feynman diagram in order to illustrate interference occurring between this background and the off-shell Higgs signal.



**Figure 3.24:** An example of a two-loop NLO contribution to the  $gg$ -channel with a closed fermion loop. The dotted black line mirrored the right Feynman diagram in order to illustrate interference occurring between this background and the off-shell Higgs signal.



**Figure 3.25:** An example of a two-fermion-loop NLO contribution to the  $gg$ -channel with a closed fermion loop. The dotted black line mirrored the right Feynman diagram in order to illustrate interference occurring between this background and the off-shell Higgs signal.



**Figure 3.26:** An example of a NLO contribution to the  $qg$ -channel with a closed fermion loop. The dotted black line mirrored the right Feynman diagram in order to illustrate interference occurring between this background and the off-shell Higgs signal.

### Massive-quark dependence

Results for off-shell Higgs production, including its quark-mass dependence at LO, are already known for both massive- and massless-quarks [87]. Additionally, NLO two-loop  $gg \rightarrow WW$  amplitudes with quark-mass dependence are available [104], but are only valid above  $m_{VV}$  as they treat the vector bosons as on-shell. As the dependence of the cross-section is expected to be significantly affected by the masses of the quarks inside the quark loops, their effect has to be accounted for.

Recently, NLO calculations for  $gg \rightarrow ZZ$  have become available in the limit  $m_t \rightarrow \infty$ , which are only valid below  $m_{t\bar{t}}$  [105–107]. These calculations are used by the  $H \rightarrow ZZ^*$  off-shell ATLAS analysis [108], applying the  $m_t$  approximation for  $ZZ$  with the addition of a large uncertainty above the  $m_{t\bar{t}}$  threshold. In the  $gg \rightarrow (H^*) \rightarrow WW$  process, the calculation is complicated through to the need to account for the mixing of massive top and bottom quarks. This makes it unable to expand around  $1/m_t$  in the  $gg \rightarrow (H^*) \rightarrow WW$  calculation.

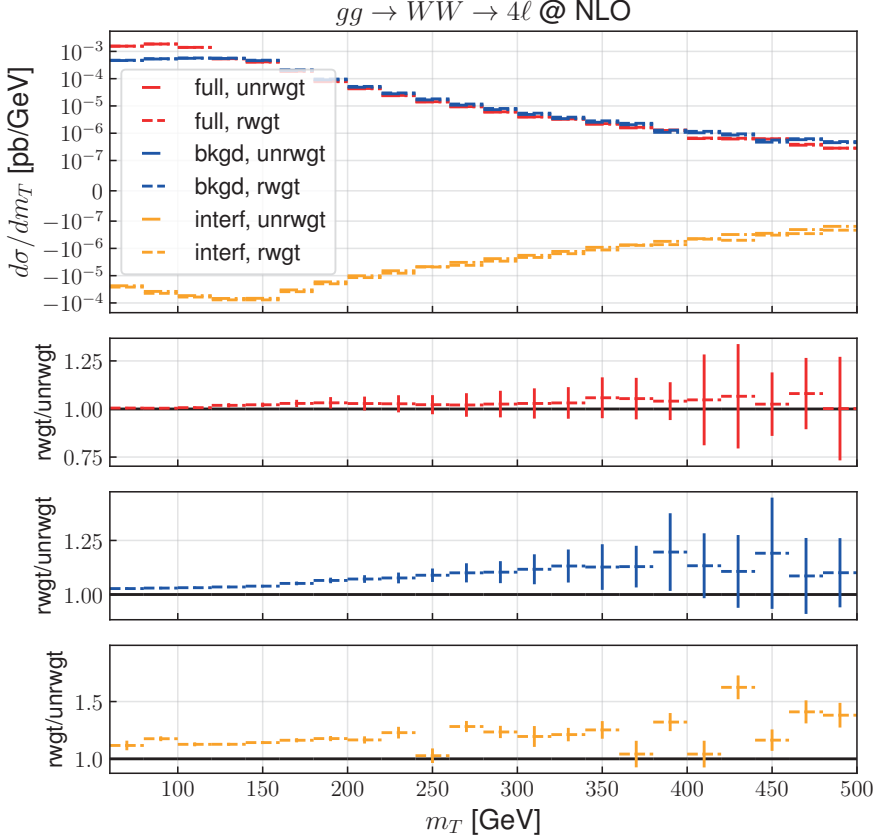
In order to remedy this, a novel procedure published in Ref. [1] is employed to account for heavy quark-mass dependence in the loops for a  $WW$  measurement. Only the four light flavour massless-quark amplitudes at the two-loop level  $\mathcal{A}_B^{\text{two-loop}, WW}(d, u, s, c)$  are computed using ggVVamp. The two-loop helicity amplitudes then get reweighed using ratios of the massive and massless quark helicity amplitudes, all known at one-loop:

$$\mathcal{A}_B^{\text{two-loop}, WW} \approx \mathcal{A}_B^{\text{two-loop}, WW}(d, u, s, c) \times \frac{\mathcal{A}_B^{\text{one-loop}, WW}(d, u, s, c, \mathbf{t}, \mathbf{b})}{\mathcal{A}_B^{\text{one-loop}, WW}(d, u, s, c)}.$$

Here,  $\mathcal{A}_B^{\text{one-loop}, WW}(d, u, s, c, \mathbf{t}, \mathbf{b})$  represents the one-loop amplitude with full mass dependence for third-generation quarks, while  $\mathcal{A}_B^{\text{one-loop}, WW}(d, u, s, c)$  denotes the one-loop helicity amplitude with four massless quarks.

Using this reweighing scheme, NLO corrections to the LO  $gg \rightarrow WW$  process are generated, accounting for the mixing of massive and massless quarks in one- and two-

loop diagrams. These corrections impact the background sample by up to +15%-20% at large transverse mass  $m_T$ ; for the interference contribution, the impact is 15% inclusively and increasing in the tail of  $m_T$ , as shown in Fig. 3.27. Therefore, it is crucial to consider these effects, even if approximated through a reweighting scheme.



**Figure 3.27:** The fixed-order NLO differential distribution of  $m_T$  for  $gg \rightarrow e\mu\nu\nu$  is shown, comparing results with and without heavy-quark-mass reweighting in the virtual amplitude. It covers the full, background (bkgd), and interference (interf) contributions. Lower panels show bin-by-bin ratios without to with reweighting. The nominal prediction uses a symlog scale [72].

### Generator settings and K-factor details

The Monte Carlo generator POWHEG-BOX v2 [70–72], built with the gg4l process, generates the corrections at fixed-order. The fixed-order K-factors are calculated separately for the signal  $gg \rightarrow H^* \rightarrow WW$ , the background  $gg \rightarrow WW$ , and the inclusive  $gg \rightarrow (H^*) \rightarrow WW$  production, including the interference.

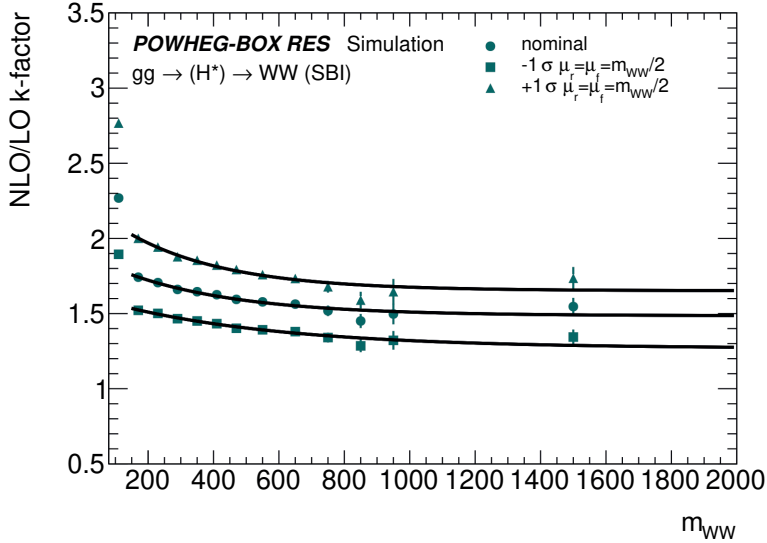
The generation uses the PDF set NNPDF3.0nnlo [93] with a strong coupling constant  $\alpha_s$  set to 0.0118. Leptons are required to have transverse momentum  $p_T > 5$  GeV, and the invariant mass  $m_{WW}$  must be greater than 1 GeV in order to avoid divergences. The  $\mu_R$  and  $\mu_F$  scale are set to  $m_{WW}/2$ , with variations including  $m_{WW}/4$  (down) and  $m_{WW}$

(up). For  $\alpha_s$ , the nominal value is 0.0118, with variations of 0.0117 (down) and 0.0119 (up). The uncertainty on the PDF is considered negligible and is estimated using LO PDF variations.

Due to the complexity of the calculations and the low cross-section in the region of interest, the K-factors exhibit some statistical fluctuations. An inclusive K-factor is obtained for the region  $m_{WW} < 140$  GeV, while a polynomial is fit to the distribution from  $m_{WW} > 140$  GeV to 2 TeV.

### Differential K-factor results and application

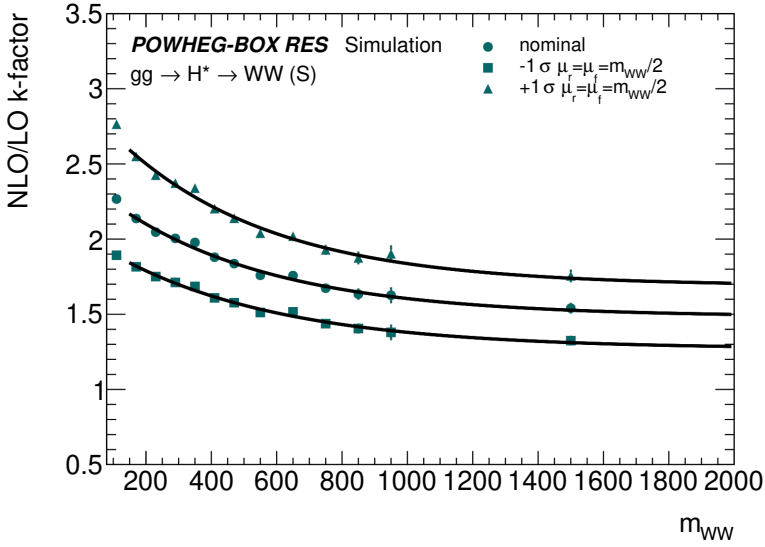
The fixed-order cross-section calculations of the differential cross-section  $d\sigma/dm_{WW}$   $\text{pb GeV}^{-1}$  are used to determine the K-factors shown in Figs. 3.28 to 3.30. The NLO uncertainties for  $\alpha_s$  are around 2% and are considered as overall normalization uncertainties.



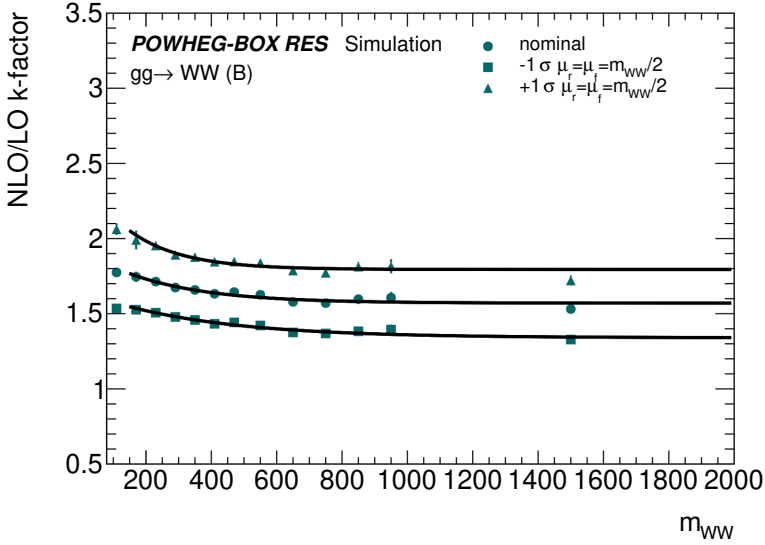
**Figure 3.28:** Fixed-order differential LO  $\rightarrow$  NLO K-factors for the  $gg \rightarrow (H^*) \rightarrow WW$  process.

Since the NLO and NNLO K-factors for the  $gg \rightarrow WW \rightarrow \ell\nu\ell\nu$  process are calculated relative to LO calculations, without any additional partons, and because the Sherpa+OpenLoopssamples include both the 0-jet and 1-jet phase-space at LO using a shower merging scheme, higher-order corrections must be applied in a step-by-step manner, in order to avoid double counting of QCD radiation. This procedure involves the usage of intermediate pure LO samples generated by MadGraph5\_aMC@NLO+Pythia8[109, 110] to achieve an accurate normalization and kinematic distributions of the  $m_{WW}$  invariant mass.

The higher-order corrections are applied to the Sherpa+OpenLoopssamples in three steps. First, higher-order QCD corrections calculated by POWHEG-BOX v2 are used to reweigh



**Figure 3.29:** Fixed-order differential LO → NLO K-factors for the  $gg \rightarrow H^* \rightarrow WW$  part of the  $gg \rightarrow (H^*) \rightarrow WW$  process.



**Figure 3.30:** Fixed-order differential LO → NLO K-factors for the  $gg \rightarrow WW$  part of the  $gg \rightarrow (H^*) \rightarrow WW$  process.

LO MadGraph5\_aMC@NLO+PythiaSBI, S and B samples as a function of  $m_{WW}^{\text{truth}}$  to NLO precision. This allows the corrected  $m_{WW}^{\text{truth}}$  distribution to serve as a reference for the differential cross-section. Next, the  $m_{WW}^{\text{truth}}$  distributions of the Sherpa+OpenLoopsLO 0+1 jet samples are reweighed to match the corrected  $m_{WW}^{\text{truth}}$  distribution from the first step. Finally, the weights as a function of  $m_{WW}^{\text{truth}}$  are applied to the  $m_{WW}^{\text{truth}}$  parton information on-the-fly during the analysis. This process reweighs the overall cross-section of the Sherpa+OpenLoopsLO 0+1 jet samples to NLO in QCD while preserving the kinematic jet information in the 1-jet region and avoids double counting.

After the application of the differential NLO K-factors, two constant higher-order K-factors are applied to the  $gg \rightarrow (H^*) \rightarrow WW$  samples: a 20% NNLO/NLO correction [95, 111, 112], and a 10% N<sup>3</sup>LO/NNLO correction [97] derived from an on-shell dominated region. Without more tailored corrections available for corrections above NLO, these were uniformly applied to the S, B, and SBI samples.

### 3.4.6 Electroweak Higgs production

The VBF process ranks second in off-shell Higgs production rate and it closely resembles the  $VH$  process, which ranks third. Both production modes form tree level diagrams for electroweak Higgs production and share identical initial and final states, causing destructive interference with the continuum background.

Similar to the  $gg \rightarrow (H^*) \rightarrow WW$  process, the cross-section as a function of  $\mu_{\text{off-shell}}$  is described as:

$$\begin{aligned} \sigma_{\text{SBI}}(\mu_{\text{off-shell}}) &\propto \mu_{\text{off-shell}}^{qq \rightarrow H \rightarrow WW + jj} \cdot \sigma_{qq \rightarrow H^* \rightarrow WW + jj \rightarrow \ell\nu\ell\nu + jj} \\ &+ \sigma_{qq \rightarrow WW + jj \rightarrow \ell\nu\ell\nu + jj} + \sqrt{\mu_{\text{off-shell}}^{qq \rightarrow H \rightarrow WW + jj}} \cdot \sigma_{\text{interference}}. \end{aligned}$$

Since the available Monte Carlo simulators cannot produce stand-alone Higgs-mediated diagrams (signal sample) for electroweak Higgs production, that include the t-channel Higgs mediated diagrams, a different strategy is adopted to compute the necessary electroweak off-shell Higgs samples.

First, a sample excluding the Higgs boson is produced, namely  $qq \rightarrow WW + jj \rightarrow \ell\nu\ell\nu + jj$  (EW B). Next, the full  $qq \rightarrow (H^*) \rightarrow WW + jj \rightarrow \ell\nu\ell\nu + jj$  process is generated (EW SBI). At last, the full process is simulated again with  $H \rightarrow VV$  couplings set to  $20^{1/4}$  times their SM strength (EW SBI20), which accounts for the coupling entering the amplitudes twice (in production and decay ( $20 \cdot \mu_{\text{off-shell}}^{\text{VBF}} \propto (20^{1/4} \cdot g_{H \rightarrow WW}^{\text{off-shell}})^4$ )), with the amplitudes squared to get the cross-section. The SBI20 sample is chosen because it sufficiently suppresses the on-shell component still present in the SBI sample and therefore limits its interference with the background during simulation.

Since the off-shell component of the Higgs signal only depends on the coupling strength to the power of 4, this modification creates a sample with  $qq \rightarrow (H^*) \rightarrow WW + jj \rightarrow \ell\nu\ell\nu + jj$  ( $\mu_{\text{off-shell}} = 20$ ). The on-shell component of the signal is adjusted back to its SM yield by increasing the Higgs width,  $\Gamma_H$ , to  $20 \times \Gamma_H^{\text{SM}}$ . This tuning of the Higgs

width is performed for the purpose of creating a sample with  $\sigma_{\text{SBI}}(\mu_{\text{off-shell}}^{\text{VBF}})$  and does not affect its  $m_{WW}$  prediction. The remaining on-shell component, generated by a  $VH$  process where a  $W$  boson (originating from the Higgs) decays hadronically, is cut out of the samples after simulation, selecting  $|m_H - 125 \text{ GeV}| > 1 \text{ GeV}$ .

The generated samples can be expressed in terms of linear combinations of the individual cross-section contributions  $\sigma_{qq \rightarrow H \rightarrow WW+jj}$ ,  $\sigma_{\text{interference}}$  and  $\sigma_{qq \rightarrow WW+jj}$ . This gives the relations

$$\begin{aligned} \text{EW SBI :} \quad & \sigma_{qq \rightarrow (H^*) \rightarrow WW+jj \rightarrow \ell\nu\ell\nu+jj}(\mu_{\text{off-shell}} = 1) = 1 \cdot \sigma_{qq \rightarrow H \rightarrow WW+jj} \\ & \quad \quad \quad \sigma_{qq \rightarrow WW+jj \rightarrow \ell\nu\ell\nu+jj} + \sqrt{1} \cdot \sigma_{\text{interference}} \\ \text{EW SBI20 :} \quad & \sigma_{qq \rightarrow (H^*) \rightarrow WW+jj \rightarrow \ell\nu\ell\nu+jj}(\mu_{\text{off-shell}} = 20) = 20 \cdot \sigma_{qq \rightarrow H \rightarrow WW+jj} \\ & \quad \quad \quad + \sigma_{qq \rightarrow WW+jj \rightarrow \ell\nu\ell\nu+jj} + \sqrt{20} \cdot \sigma_{\text{interference}} \\ \text{EW B :} \quad & \sigma_{qq \rightarrow WW+jj \rightarrow \ell\nu\ell\nu+jj}(\mu_{\text{off-shell}} = 0). \end{aligned}$$

These relations can be summarized in the following matrix equation

$$\begin{bmatrix} \text{EW SBI} \\ \text{EW SBI20} \\ \text{EW B} \end{bmatrix} = \begin{bmatrix} 1 & 1 & 1 \\ 20 & \sqrt{20} & 1 \\ 0 & 0 & 1 \end{bmatrix} \cdot \begin{bmatrix} \sigma_{qq \rightarrow H \rightarrow WW+jj} \\ \sigma_{\text{interference}} \\ \sigma_{qq \rightarrow WW+jj} \end{bmatrix}.$$

Using the formula shown below, any  $\mu_{\text{off-shell}}^{\text{VBF}}$  hypothesis can be produced from the SBI, SBI20 and B samples, of which [Tab. 3.1](#) shows the cross-section information.

$$\begin{aligned} \sigma_{\text{SBI}} = & \frac{\mu_{\text{off-shell}} - \sqrt{\mu_{\text{off-shell}}}}{20 - \sqrt{20}} \cdot \sigma_{\text{SBI20}} + \frac{20\sqrt{\mu_{\text{off-shell}}} - \sqrt{20}\mu_{\text{off-shell}}}{20 - \sqrt{20}} \cdot \sigma_{\text{SBI1}} \\ & + \frac{(\sqrt{\mu_{\text{off-shell}}} - 1)(\sqrt{\mu_{\text{off-shell}}} - \sqrt{20})}{\sqrt{20}} \cdot \sigma_{\text{B}} \end{aligned}$$

### Electroweak Higgs simulation

The VBF and  $VH$  production modes are generated using MadGraph5\_aMC@NLO and PYTHIA 8 with the NNPDF3.0nlo [93] set and the A14 tune [113] for parton showering and hadronization. The tendency of the default Pythia settings to produce excessive radiation in the central region ( $|\eta| < 2.5$ ) is mitigated using the dipole-recoil scheme [114], which ensures that only a single final-state parton recoils from each emission.

The VBF samples are generated with MadGraph5\_aMC@NLO at LO in QED, with the number of QCD vertices required to be zero, in order to produce  $WW$  production in association with two jets. The full vector boson decay is present in the matrix element generated by MadGraph5\_aMC@NLO and accounts for off-shell effects of the intermediate vector bosons, as well as the Higgs boson.





## THE STANDARD MODEL AS AN EFFECTIVE FIELD THEORY

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In the absence of direct evidence for new physics, efforts focus on systematically searching for deviations from the SM using the available data. This chapter introduces a more systematic, model-independent strategy for the interpretations of BSM effects. This chapter explains the SM-EFT [115] and its application in analyses.

### 4.1 Introduction to effective field theory

Consider an ultraviolet (UV) complete QFT, able to describe the kinematics of new physics present at an energy scale  $M$ . In this theory, the parameter  $M$  is the energy scale of the high-mass particle, the new physics energy scale, and  $m$  is the energy scale of the known low-energy phenomena. The couplings between the particles at energy scale  $M$  and  $m$  yield a perturbative expansion in the limit  $M \gg m$  with terms of the form  $\alpha^r \log^s(m/M)$  [116] as a function of their coupling strength  $\alpha$ .

Because of this expansion, the interactions of both types of physics can be decoupled due to the different scales at which they reside [117]. However, in certain phase-spaces and depending on the absolute separation of energy scales, the coupling terms become large and the perturbative expansion breaks down; a UV-complete QFT would describe the new physics at scale  $M$  and the low-energy physics at scale  $m$ . However, since it is a perturbative expansion, it diverges due to large logarithmic terms.

A solution to these IR divergences is to use the concept of effective field theory (EFT). This absorbs the diverging terms into a cutoff scale  $\Lambda$ , where  $\Lambda$  is the scale of new physics. A matching procedure trades the IR divergences of the UV-complete theory to the UV divergences of the EFT. The Wilson coefficients derived during the matching procedure, at matching scale  $\Lambda$ , run from the UV to the IR scale using the renormalization group equations in order to alleviate the remaining UV divergences [116]. The EFT can now describe phenomena caused by physics at energy scale  $M$  on cutoff scale  $\Lambda$  by absorbing the effects of new physics at higher energies into the parameters of lower scale physics, which are studied by experiments operated at energy scale  $m$ .

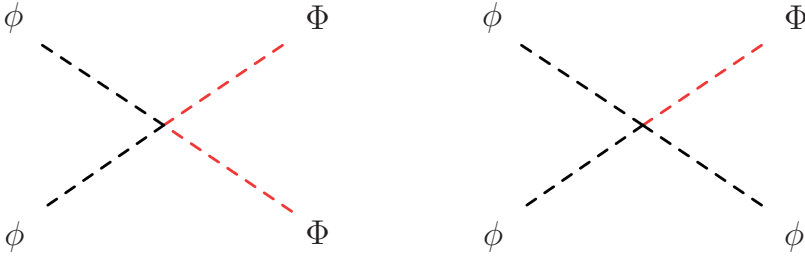
The EFT can be derived in the top-down or bottom-up approach [118]. While top-down EFTs make certain assumptions about high-energy phenomena to specify or justify the low-energy theory, bottom-up EFTs are developed to parameterize unknown observations into known parameters or, in the case of this thesis, aid in searches for currently unknown

high-energy physics. The use of a scalar toy model [116], can illustrate the concept of EFT in a simple way and demonstrate how an EFT can be used to describe low-energy phenomena of a high-energy theory.

The toy model is a UV-complete theory and a scalar field theory of  $\phi$  and  $\Phi$  with a mass of  $m$  and  $M$  respectively. Consider the terms

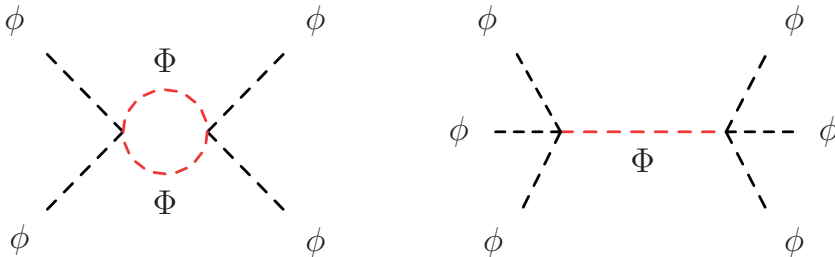
$$\begin{aligned}\mathcal{L}_{\text{kinetic+mass terms}}^{\text{UV-complete}} &= \frac{1}{2}(\partial_\mu \phi)^2 - \frac{1}{2}m^2\phi^2 + \frac{1}{2}(\partial_\mu \Phi)^2 - \frac{1}{2}M^2\Phi^2 \\ \mathcal{L}_{\text{interaction terms}}^{\text{UV-complete}} &= -\frac{1}{6}a\phi^3 - \frac{1}{2}b\phi^2\Phi - \frac{1}{4}\kappa\phi^2\Phi^2 - \frac{1}{6}\rho\phi^3\Phi - \frac{1}{24}\eta\phi^4.\end{aligned}$$

In this lagrangian, the constants  $a, b, \kappa, \rho, \eta$  are dimensionless and the fields  $\phi$  and  $\Phi$  are real scalar fields. The two interaction terms  $-\frac{1}{4}\kappa\phi^2\Phi^2$  and  $-\frac{1}{6}\rho\phi^3\Phi$  present in the lagrangian  $\mathcal{L}^{\text{UV-complete}}$  are the ones that are of interest to us. These are the two four-point  $2 \rightarrow 2$  and  $3 \rightarrow 1$  interactions, shown in Fig. 4.1.



**Figure 4.1:** The Feynman rules of the two interaction terms  $-\frac{1}{4}\kappa\phi^2\Phi^2$  and  $-\frac{1}{6}\rho\phi^3\Phi$ .

This information is used in the UV-complete theory and matched to an EFT featuring only  $\phi$  fields. By using the Feynman rules depicted in Fig. 4.1 to go to the IR theory, where only the low mass particles  $\phi$  are produced on-shell, the diagrams in Fig. 4.2 are constructed.



**Figure 4.2:** The Feynman diagrams which can be constructed using the two interaction terms  $-\frac{1}{4}\kappa\phi^2\Phi^2$  and  $-\frac{1}{6}\rho\phi^3\Phi$ .

In these loops, the field  $\Phi$  appears only virtually and its final states are the only ones that can be observed at low energies, since the  $\Phi$  field cannot be directly produced. These

processes are described by the EFT lagrangian

$$\mathcal{L}^{\text{EFT}} = \frac{1}{2}(\partial_\mu \phi)^2 - \frac{1}{2}m^2\phi^2 - \frac{1}{4!}C_4\phi^4 - \frac{1}{6!}\frac{1}{M^4}C_2\phi^6 - \frac{1}{8!}\frac{1}{M^4}C_8\phi^8 + \dots$$

The UV theory parameters  $\kappa$  and  $\rho$  are now matched onto the EFT parameters  $C_k$  using dedicated matching calculations that have to be performed separately for each matched UV theory. This method effectively integrates out [116] the heavy particle  $\Phi$  and expresses its influence on low-energy physics as an expression in local interactions of the light fields  $\phi$ , suppressed by powers of the heavy mass scale  $M$ .

## 4.2 The Standard Model Effective Field Theory

The SM is currently the best available description of fundamental interactions at the microscopic scale. However, because of a number of shortcomings described in § 2.7, we can think of the SM as an EFT of a more fundamental theory with new degrees of freedom at energy scale  $\Lambda$ . The operator basis in which the SM-EFT is to be expressed can be chosen relatively freely, although certain choices in basis are made to simplify the calculations.

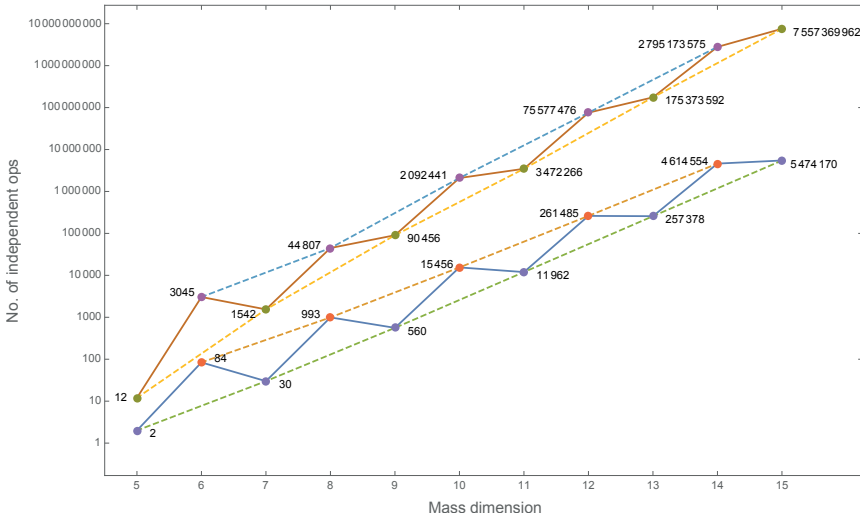
The derivation of the SM-EFT basis and the matching calculations of the SM towards this EFT are performed in [119–121] and constitute a complete basis of operators up to mass dimension (D) = 6, which are suppressed by  $\Lambda^{D-4}$ , such that the overall mass dimension is 4. The mass dimension-4 is required such that when integrating the lagrangian over the dimension-4 space, the resulting action  $S = \int \mathcal{L} dx^4$  is dimensionless. In the SM-EFT basis, the  $SU(3)_c \times SU(2)_L \times U(1)_Y$  symmetries are respected and its input variables are taken from the SM. This yields the following lagrangian for the SM-EFT

$$\begin{aligned} \mathcal{L}_{\text{SM-EFT}} = & \mathcal{L}_{\text{SM}}^{(4)} && \text{(Mass dimension D = 4)} \\ & + \sum_i \frac{c_i^{(5)}}{\Lambda} \mathcal{O}_i^{(5)} + \sum_k \frac{c_k^{(7)}}{\Lambda^3} \mathcal{O}_k^{(7)} && \text{(Odd Mass dimension D = 5,7)} \\ & + \sum_j \frac{c_j^{(6)}}{\Lambda^2} \mathcal{O}_j^{(6)} + \sum_l \frac{c_l^{(8)}}{\Lambda^4} \mathcal{O}_l^{(8)} && \text{(Even Mass dimension D = 6,8)} \\ & + \sum_m \frac{c_m^{(D)}}{\Lambda^{D-4}} \mathcal{O}_m^{(D)} + \dots && \text{(Higher Mass dimension D > 8).} \end{aligned}$$

Here, the  $\mathcal{L}_{\text{SM}}^{(4)}$  is the SM lagrangian, the  $\mathcal{O}_i^{(5)}$  and  $\mathcal{O}_k^{(7)}$  are the odd mass dimension operators, the  $\mathcal{O}_j^{(6)}$  and  $\mathcal{O}_l^{(8)}$  are the even mass dimension operators and the  $\mathcal{O}_m^{(D)}$  are higher dimension operators. The  $c_{\alpha=(i,j,k,l,m)}^D$  are the Wilson coefficients corresponding with their associated mass dimension. The complete basis of operators in each mass dimension has to be known in order to make the EFT, up to that mass dimension, renormalizable.

In the SM-EFT, the mass dimension-5 operator type is known as the Weinberg operator [122], which is the only dimension-5 operator that can be constructed from the SM fields. This operator allows for the generation of Majorana neutrino masses and the violation of leptonic number. As the neutrino masses are expected to be  $m_\nu < 1$  eV [123], the energy scales of these operators are beyond the energy reach of the LHC and therefore neglected in this work.

The leading contribution of SM-EFT operators at the LHC, assuming a scale of  $\Lambda = 1$  TeV, are the dimension-6 operators. The full number of operators from mass dimension-6 is shown in Fig. 4.3, along with the higher mass dimensions. When assuming baryon and lepton number conservation, the number of operators is further reduced to an orthogonal basis of 2499 operators. When assumptions are made on the flavour awareness of the new physics, i.e., making it flavour agnostic, it can be further reduced to a manageable set of 59 orthogonal operators. Under certain flavour symmetries, the operators are agnostic about the flavour of the fields they are constructed from. In both analysis described in Chapter 13 and Chapter 12 the  $U(3)^5$  flavour scheme is applied, given by the symmetry group  $U(3)_q \times U(3)_u \times U(3)_d \times U(3)_l \times U(3)_e$  [124, 125].



**Figure 4.3:** Growth of the number of independent operators in the SM-EFT up to mass dimension 15. Points joined by the lower solid line when no flavour distinction is being made; those joined by the upper solid line are for three generations treated independently. Dashed lines are to guide the eye to the growth of the even and odd mass dimension operators in both cases [126].

The dimension-7 operators are suppressed by a factor of  $\Lambda^3$  and are expected to be less relevant than the dimension-6 operators. On top of the  $\Lambda^3$  suppression, they also violate lepton and baryon number conservation. Since lepton and baryon number violation is not observed at the LHC, the terms are expected to be small and therefore neglected.

At last, the dimension-8 operators are suppressed by  $\Lambda^4$  and are therefore expected to

be even less relevant than the dimension-7 operators. However, despite the fact they are subleading in the EFT expansion, this does not mean they can be safely neglected in all cases. As their interference with SM operators could lead to corrections comparable in size to the cross-section components of dimension-6 terms, for example processes containing triple gauge couplings [127].

### 4.3 The Warsaw basis

The choice of basis is not unique; however, the field, encompassing both theorists and experimentalists, has converged on adopting the Warsaw basis [128]. This basis has recently been implemented in Monte Carlo generators which can be used to make predictions for signatures in the LHC.

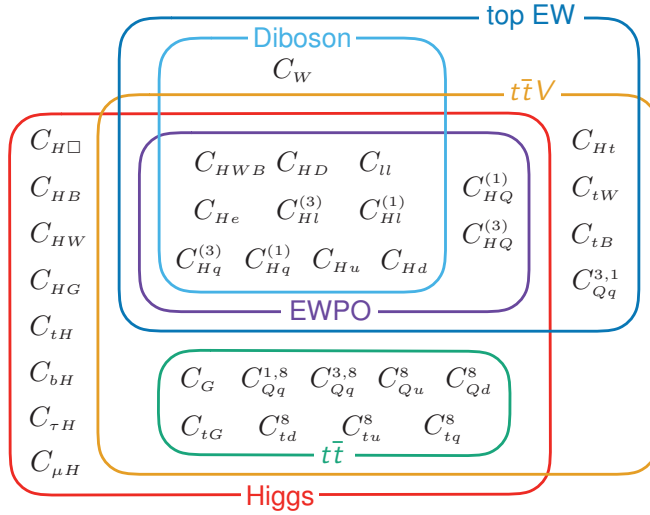
One of the most common models is the SMEFTsim package [124, 129, 130]. It can simulate the full set of dimension-6 operators at leading order in QED and QCD using no flavour assumptions, minimal flavour violating or maximal flavour symmetry labelled as the  $U(3)^5$  scheme. Following this model, the SMEFT@NLO [125, 131] model was published, which allows for the simulation of dimension-6 operators at the one-loop level in QCD in the topU3l flavour scheme. This flavour scheme ensures the third generation of quarks is modelled independently from the first two generations.

### 4.4 Measurements of operators and new physics

The advantage of the SM-EFT approach is that it provides a consistent model to describe the effect of new physics across all measurable physics processes at the LHC and thus allows to combine such measurements in a coherent way. Once combinations are performed, the measurements can exploit deviations from the SM and find patterns in operator limits.

The sharing of dimension-6 operators across the different particle physics sectors of top,  $t\bar{t}V$ , Higgs, diboson and electroweak precision observables (EWPOs), illustrated in Fig. 4.4, can be taken as an example. In this case the different classes of physics processes at the LHC contribute to jointly measure the various SM-EFT operators. The figure also illustrates that individual analyses are often sensitive only to a subset or linear combination of operators. Consequently, combining multiple measurements is necessary to fully explore the dimension-6 operator space.

When limits of operators are measured at the LHC, they are published as limits on their Wilson coefficients using a particular value of  $\Lambda$ . Since the SM-EFT is expanded around  $c_i/\Lambda$ , experimental constraints on EFT operators should be interpreted in terms of this ratio. Once these limits are supplied, they can be mapped back onto parameters of any UV-complete model. This SM-EFT approach ensures reusability of constraints acquired from high-energy physics measurements. Additionally, when a deviation from the SM is found and interpreted in terms of EFT operators, theorists are able to directly map constraints on Wilson coefficients to their newly developed UV models.



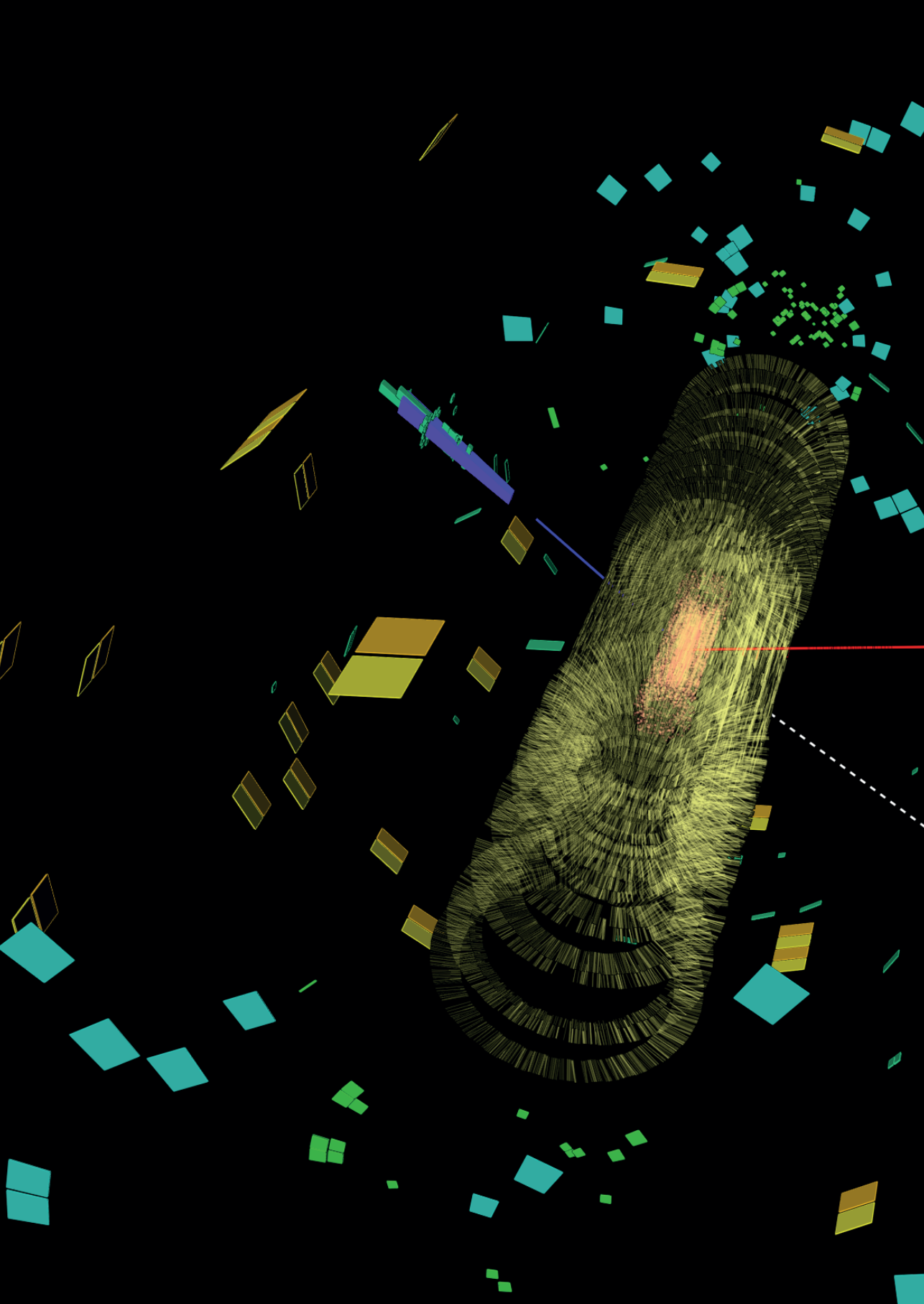
**Figure 4.4:** An illustration of the sharing of dimension-6 operators across different the particle physics sectors. The figures shows the top,  $t\bar{t}V$ , Higgs and diboson sectors, as well as measurements of EWPOs. Examples of EWPOs are  $\Gamma_Z$  and  $m_W$ . Taken from Ref. [132].

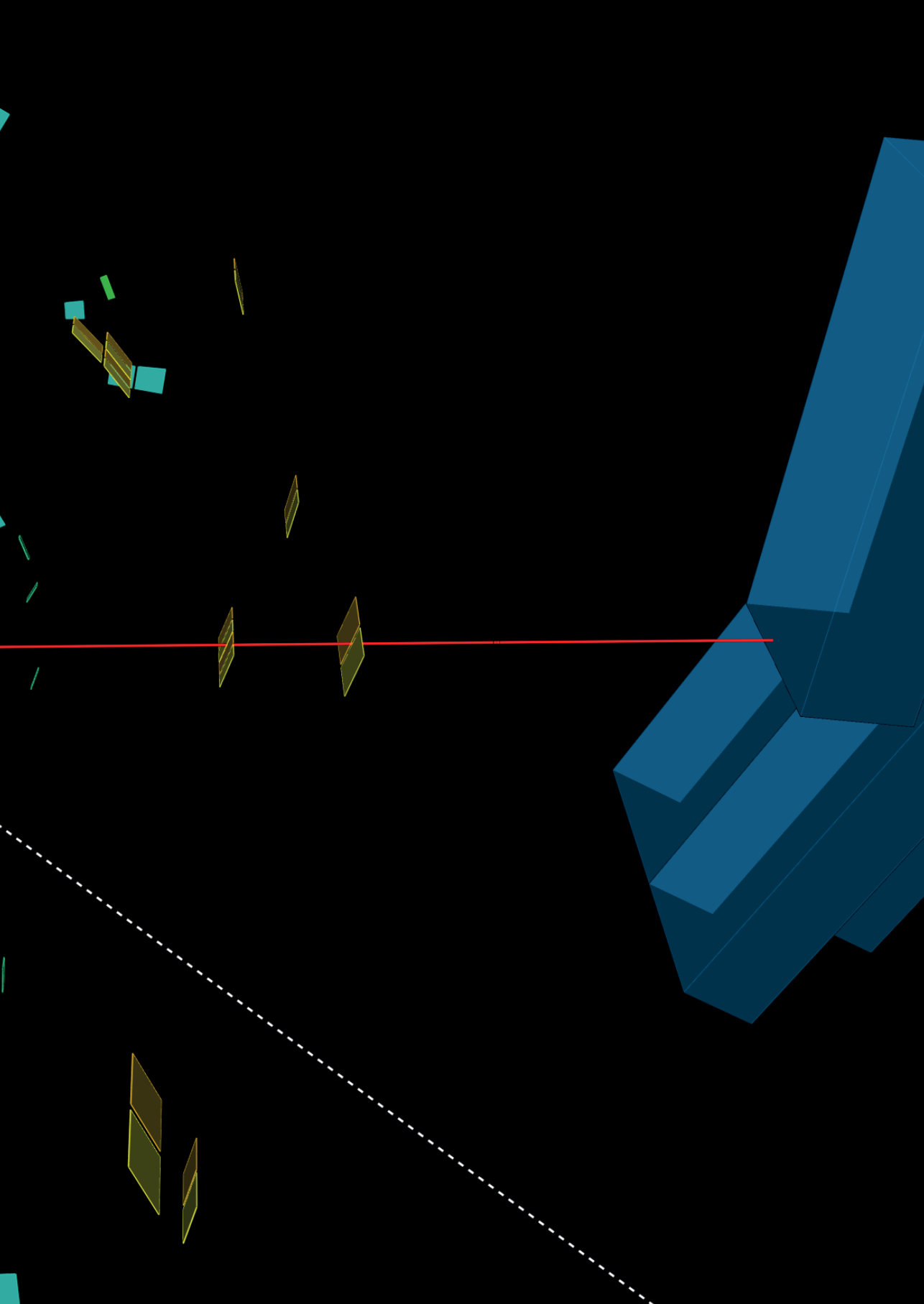
The SM-EFT is a powerful tool to systematically and efficiently search for new physics at the LHC and beyond. However, one should note that the SM-EFT is a largely model-independent approach, but that it still makes some assumptions about the nature of new physics. The SM-EFT assumes that the new physics is weakly coupled, assumes a SM like structure, and assumes that the new physics is heavy enough to be integrated out. Due to these assumptions, Higgs width deviations through decays into low-mass particles are not captured. As a result, SM-EFT searches and the Higgs width measurements described in Chapter 11 complement each other, offering a more holistic approach to identify potential deviations from the SM and uncover new physics.

PART II

EXPERIMENT









## THE LARGE HADRON COLLIDER

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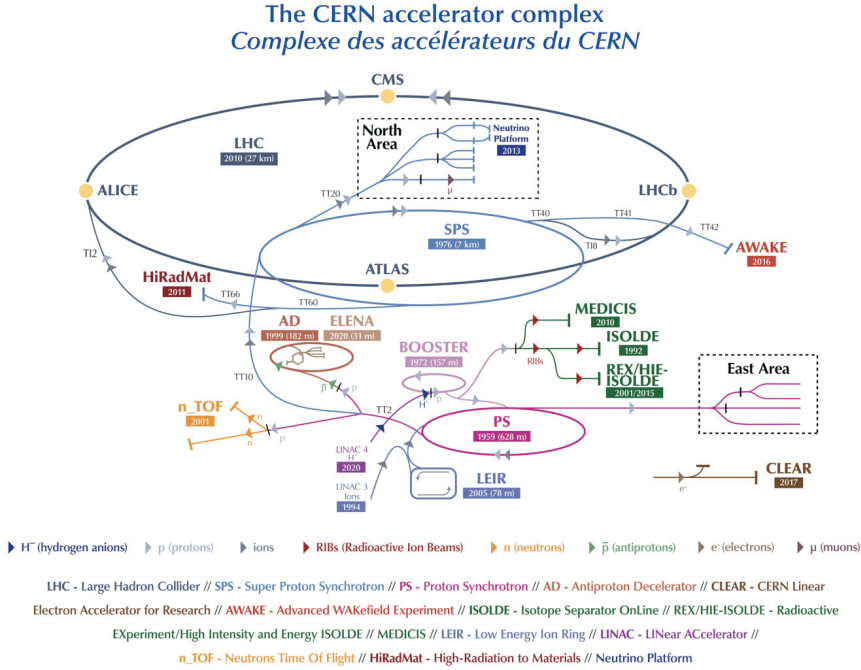
The results in this thesis would not have been possible without the successful operation of not only the LHC, but also the various pre-accelerators distributed around the CERN complex. These accelerators are able to produce a high number of Higgs bosons in a controlled experimental environment. Currently, this is the only accelerator complex in the world where Higgs bosons are produced with such abundance so they can be studied with high precision. In this chapter the LHC and the CERN accelerator complex are introduced.

### The Accelerator

The accelerator which produces the data for this thesis is the LHC. This circular collider is partially located under the canton of Geneva and is part of the CERN laboratory. The collider is housed inside the 27 km long tunnel located 100 m underground which previously hosted the Large Electron Positron Collider (LEP). The LHC has been operating successfully since 2010 and was designed to collide protons at a centre-of-mass energy of 14 TeV.

Run-1 activity refers to the years of data taking between 2010 and 2012. In this period the centre-of-mass energy was ramped up from  $\sqrt{s} = 7$  TeV towards  $\sqrt{s} = 8$  TeV at the end of the run. After the scheduled technical stop of two years allowing for the Phase-I upgrades to take place, Run-2 started. This run period lasted from 2015 to 2018, which also featured an increase in centre-of-mass energy towards  $\sqrt{s} = 13$  TeV. The Run-2 dataset is used in the main work of this thesis. After Run-2, Run-3 started in 2022; in this run, the centre-of-mass energy was increased further towards  $\sqrt{s} = 13.6$  TeV.

The proton's journey from pressurized hydrogen barrel all the way to  $pp$ -collisions at the interaction points of the various LHC experiments goes as follows. The protons are accelerated in a series of accelerators before they reach the LHC; the CERN accelerator complex during the start of Run-3 is depicted in [Fig. 5.1](#). First, hydrogen atoms are stripped from their electrons by a strong magnetic field to produce the protons. The protons are then accelerated in the Linac accelerator, the Proton Synchrotron Booster, the Proton Synchrotron and finally the Super Proton Synchrotron (SPS). Protons arriving in the LHC are accelerated up to about 450 GeV by the SPS. The SPS provided the beams for the UA1 and UA2 experiments where the  $W$  and  $Z$  bosons are discovered [31]; now it serves as a pre-accelerator for the LHC.



**Figure 5.1:** The accelerator complex of CERN during the start of Run-3. Taken from Ref. [133].

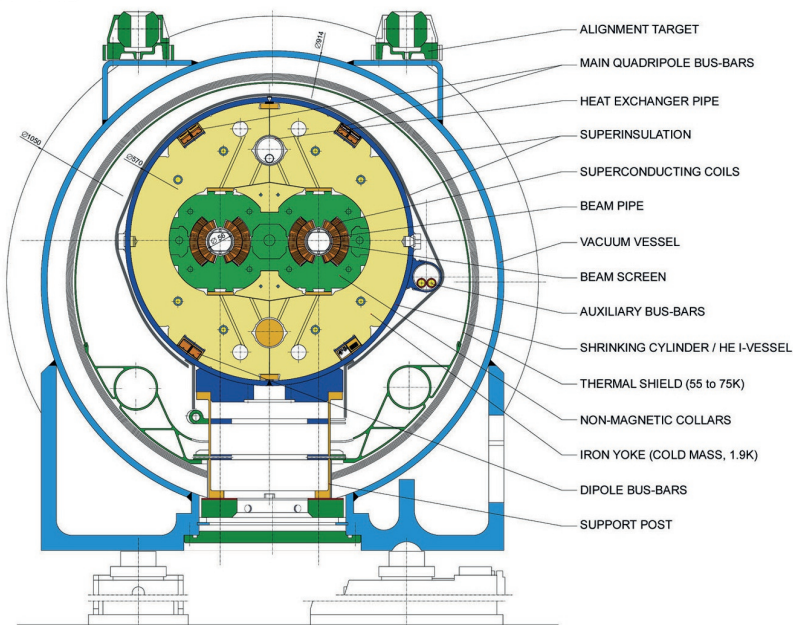
When arriving at the LHC, the protons are injected into the accelerator in bunches, in order to maximize the collision chance per bunch crossing. The design value for the number of bunches was 2808 for 1 beam, with about  $1.15 \times 10^{11}$  of protons per bunch and a bunch spacing of 25 ns. When the maximum number of bunches are circulating in the LHC beams they are considered as fully populated; the circular beam paths will then be directed to form collisions points at each of the four LHC experiments. These experiments are ALICE, CMS, LHCb and ATLAS, the detector used in this thesis.

In order to reach the desired LHC Run-2 centre-of-mass energy of 6.5 TeV per beam, the protons are accelerated using eight superconducting radio frequency cavities. These accelerators use a varying electrical field to transfer the energy stored inside the cavity of the system into the particle's momentum and drive them forward though the beam pipe. The direction of flight of the bunches is then curved, along the circular beam pipe, using superconducting niobium-titanium (NbTi) dipole magnets. Since the protons all carry positive charge, they repel each other, which increases the width and height of the bunches inside the beam pipe. The width and height of the bunches is therefore controlled using superconducting NbTi quadrupole magnets distributed around the beam pipe.

Superconductivity is the phenomenon where the electrical resistance of a material drops to zero when its temperature, its magnetic field and current are lowered below a critical

value [134]. This allows for magnets to be operated at a high current without any energy loss. So, a very high current density can be achieved. This in turn leads to the generation of strong magnetic fields, which are needed to mould and bend the proton beams in the LHC.

The magnets are kept superconductive using superfluid helium, operating at a temperature of 1.9 K. The dipole magnets [135], illustrated in Fig. 5.2, are designed to allow for a nominal current operation of about  $11.6 \times 10^3$  A, producing a peak central field of 8.3 T, with the possibility to reach 9 T. The quadrupoles require a lower peak field of about 7.5 T. Both superconducting magnets are built using copper stabilized NbTi Rutherford cables. The copper ensures a high thermal conductivity between the cooling agent and the superconducting wire, as well as a secondary current path and heat sink should the magnet quench.



**Figure 5.2:** A schematic cross-section of an LHC dipole magnet highlights the components responsible for generating the magnetic field, along with its supporting structures. Taken from Ref. [135].

A magnet quench [136] is the sudden loss of superconductivity in the magnet, which leads to a rapid increase in temperature. This temperature increase causes further loss of superconductivity, as well as evaporation of the cooling agent and can thus start a cascade effect causing all the energy stored in the magnetic field to be released on the magnet material in the form of heat. Quenches are prevented by the use of a quench protection system which dumps the energy stored in the magnet into a resistor bank or the cold yoke of the magnet itself.

The LHC luminosity ( $L$ ) can be defined as

$$L = \frac{n_b N_p^2 f_{\text{rev}}}{4\pi\sigma_T^2} F.$$

In this equation,  $N_p$  stands for the number of particles per bunch, while the number of bunches per beam is  $n_b$  and the total transverse bunch size is defined as  $\sigma_T = \sigma_x \cdot \sigma_y$ . The revolution frequency is  $f_{\text{rev}}$ , which amounts to 40 MHz at the LHC. The parameter  $F$  takes into account the luminosity reduction due to the crossing angle of the beams. In Run-2, the peak luminosity of the LHC reached  $1.9 \times 10^{34} \text{ cm}^{-2} \text{ s}^{-1}$  as the number of colliding bunch pairs and number of protons per bunch increased during the course of each year [137].

The LHC is operated in order to discover rare processes, which typically have low probability of occurrence. Therefore, a high luminosity is preferred to discover these rare processes. The production rate of an arbitrary rare event is defined as

$$N = \sigma \cdot \mathcal{L}.$$

Here,  $\sigma$  is the cross-section of the process and  $\mathcal{L}$  is the integrated luminosity of the collider defined as

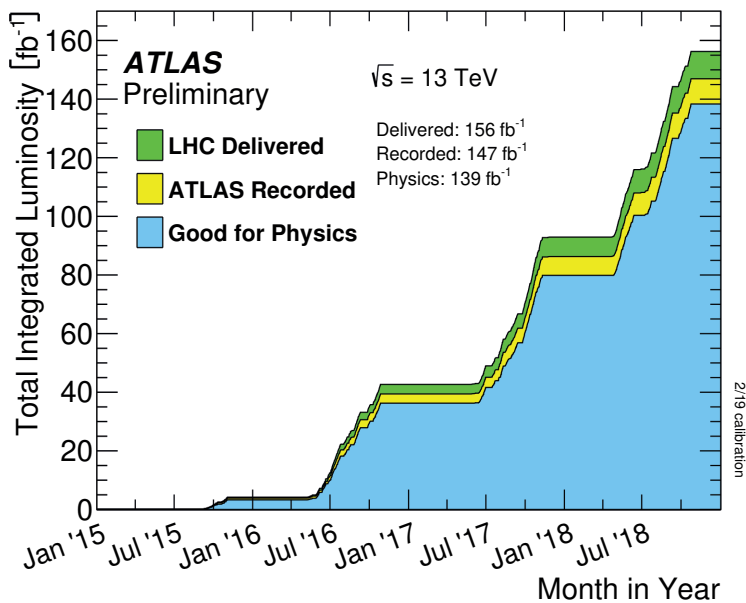
$$\mathcal{L} = \int L \cdot dt.$$

The integrated luminosity is used as a measure of the amount of events the experiment has collected and is measured using the unit of inverse barn ( $b^{-1} = 10^{24} \text{ cm}^{-2}$ ). During Run-2,  $139 \text{ fb}^{-1}$  of data passed all good run checks and was used for analysis, as shown in Fig. 5.3.

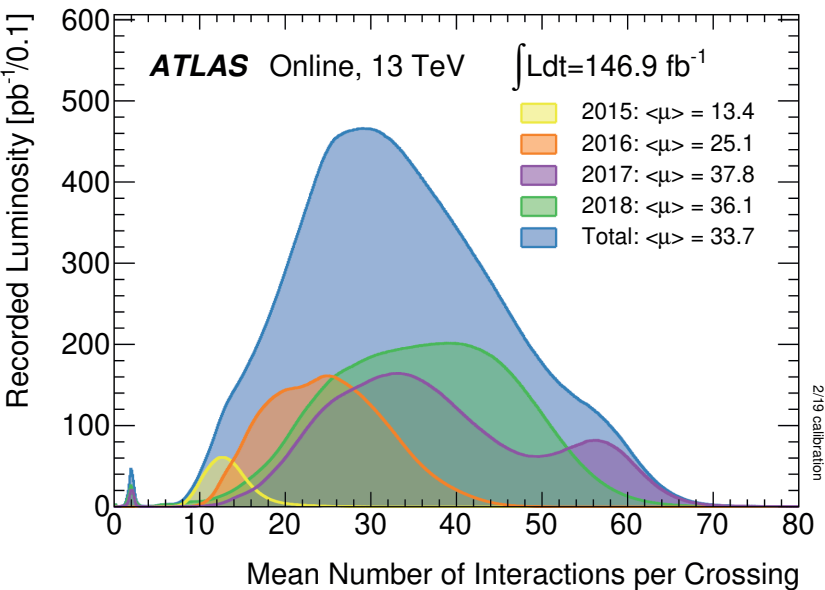
Since each bunch contains many protons, in addition to the  $pp$ -collision with the highest momentum transfer that triggers the events, several other less energetic  $pp$ -collisions also take place almost simultaneously. These collisions are referred to as pile-up collisions. When a bunch crossing occurs, the number of proton-proton interactions can be defined using the following equation

$$\mu = \frac{L\sigma_{\text{inelastic}}}{n_b f_{\text{rev}}}.$$

The amount of pile-up events decreases with the mean intensity as the beams collide. In this equation  $\sigma_{\text{inelastic}}$  is the total inelastic cross-section,  $n_b$  the number of bunches and  $f_{\text{rev}}$  the collision frequency. The distribution of the mean number of interactions of all bunch crossings as a function of the run year is shown in Fig. 5.4, and also shows the total delivered luminosity as a function of the mean number of interactions per crossing  $\langle\mu\rangle$  for the total Run-2 dataset and its constituent years at  $\sqrt{s} = 13 \text{ TeV}$ .



**Figure 5.3:** Cumulative luminosity versus time, delivered to ATLAS, recorded by ATLAS, certified to match the quality criteria. Taken from Ref. [138].



**Figure 5.4:** The recorded luminosity as a function of the mean number of interactions per crossing. Taken from Ref. [138].





## THE ATLAS DETECTOR

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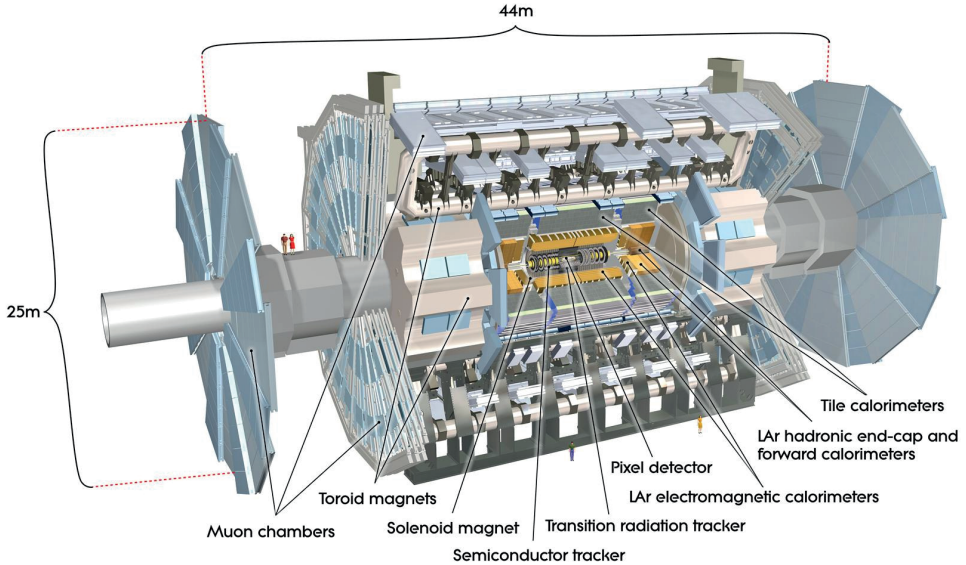
The collisions produced by the LHC and the CERN accelerator complex are reconstructed using the ATLAS detector. In this section, the ATLAS detector and its various subdetectors are explained. In particular, the detectors involved in the reconstruction the final state objects of interest in this thesis are discussed. These are mainly electrons, muons, particle jets and missing transverse momentum.

### 6.1 An overview of the ATLAS Detector

The ATLAS (A Toroidal LHC Apparatus) detector [139] is a general-purpose detector designed to detect and reconstruct a wide variety of collision processes and decay products. The LHC's operation requires ATLAS to maintain strong particle identification performance in an environment where protons collide at  $\sqrt{s} = 13$  TeV, resulting in simultaneous interactions and a high particle density.

ATLAS and its subdetectors are schematically drawn in Fig. 6.1. In order to allow the maximal number of particles to be detected by one of the ATLAS subdetectors, the detector forms a nearly closed structure. Overall, the ATLAS detector consists of a cylindrical barrel region and two endcaps, forming a closed cylinder sealed at both ends while leaving the beam pipe open. This design provides a near  $4\pi$  field of view coverage of the interaction point, maximizes the detector acceptance and allows for the reconstruction of final states as complete as possible. In the barrel region, several subdetectors are situated, layered around the interaction point to provide direct access to the particles created in the collisions. In the endcaps, layered subdetectors are present to cover the forward solid angle of the interaction point.

In order to perform the particle tracking, the Inner Detector (ID) is closest to the interaction point. The Pixel Detector (PD), Semiconductor Tracker (SCT) and Transition Region Radiation Tracker (TRT) inside the ID are immersed in a 2 T axial magnetic field provided by a superconducting solenoid magnet. The magnetic field bends the charged particle tracks, which are used to extract the momentum. The ID is surrounded by Liquid Argon (LAr) electromagnetic (EM) calorimeters both in the barrel region and the endcaps. These EM calorimeters are in turn surrounded by the hadronic calorimeters, which provide energy measurements for particles showering due to the electromagnetic force or the strong force.



**Figure 6.1:** Schematic illustration of the ATLAS detector. On the left beam pipe two people are illustrated to give sense to the sheer scale of the experiment. Taken from Ref. [140].

The subdetector furthest away from the interaction point is the Muon Spectrometer (MS). Because highly energetic muons are able to pass relatively undetected through the previously mentioned subdetectors because of their low interaction cross-section, the performance of this subdetector remains unaffected. This subdetector relies on dedicated muon tracking detectors and superconducting toroid magnets to provide a toroidal magnetic field in the barrel and endcap regions.

The ATLAS magnet system [141] consists of four superconducting magnets. The Central Solenoid (CS), the Barrel Toroid (BT) and two endcap Toroids (ECT). As each of the magnets is superconducting, dedicated cooling circuits and individual cryostats for thermal isolation are present. In order to withstand Lorentz forces and the weight of the coils, a dedicated mechanical structure is present, enclosing each coil and the structure between them. The peak field loads, which are the magnetic fields in the superconductor windings, correspond to 3.9 T, 4.1 T, 2.6 T for the BT, ECT and CS, respectively. The superconducting conductor used is NbTi, stabilized with copper and aluminium. The NbTi/Cu/Al conductor is used as a single layer coil, wound internally using a support cylinder.

The toroids are wound into flat racetrack coils using a pancake winding technique. Because of the varying sizes and stabilizer material inside the conductors used for the various magnets, their critical currents ( $I_c$ ) vary as well. The designed critical current for the conductors used in the BT, ECT and CS are 47.5 kA, 55 kA and 20.4 kA at a magnetic

field strength of 5 T and temperature of 4.2 K, respectively. The nominal operation point ( $I/I_c$ ) of the magnets lies between 40% and 20%. To maintain their superconductive state, the magnets must remain below their critical temperature ( $T_c$ ). This is achieved through a dedicated indirect cooling method, using forced-flow helium at 4.5 K, circulating through tubes welded onto the outer casing of the conductor windings.

The ATLAS detector coordinate system is defined as follows: the  $z$ -axis is set along the beam direction and the  $x - y$  plane is transverse to the beam direction. The  $x$ -axis points toward the middle of the LHC ring and the  $y$ -axis points upwards. However, when referring to locations of particles in the detector, a cylindrical coordinate system is used because of the cylindrical layout of the detector. The polar angle,  $\theta \in [0, \pi]$ , is the angle from the beam axis. The azimuthal angle,  $\phi \in [-\pi, \pi]$ , is the angle around the beam axis. The pseudo-rapidity ( $\eta$ ) is used to describe the polar angle of a particle relative to the beam axis. The rapidity ( $y$ ) and  $\eta$  coincide for massless particles and are defined as

$$y = \frac{1}{2} \ln \left[ \frac{E + p_z}{E - p_z} \right], \quad \eta = -\ln \tan \left( \frac{\theta}{2} \right).$$

The solid angle in the pseudorapidity-polar angle space is defined as  $\Delta R = \sqrt{\Delta\eta^2 + \Delta\phi^2}$  in the transverse plane. The transverse momentum ( $p_T$ ) and transverse energy ( $E_T$ ) are defined as

$$p_T = p \cdot \sin\theta, \quad E_T = E \cdot \sin\theta.$$

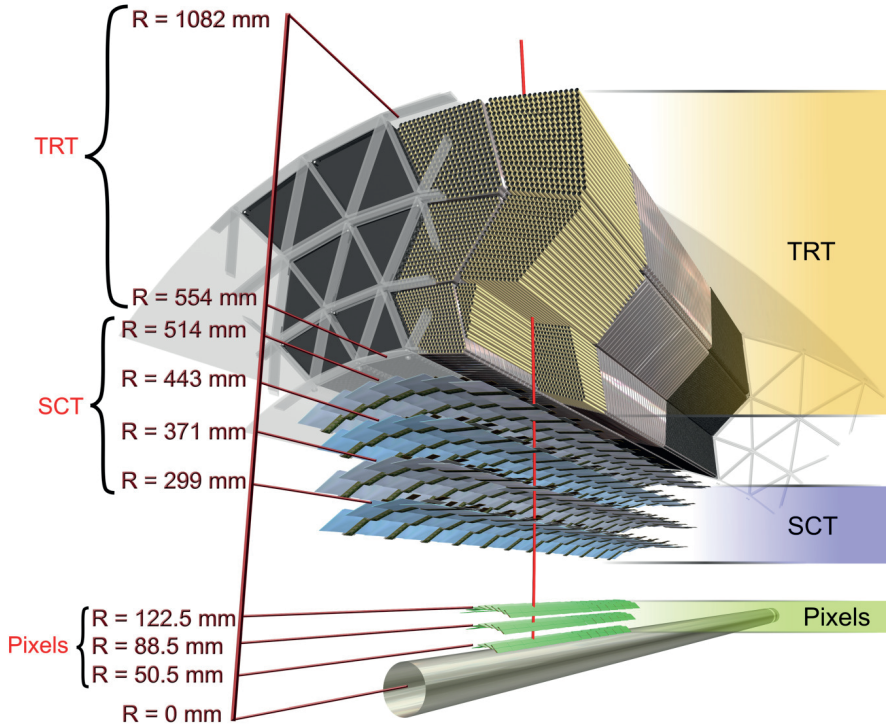
## 6.2 The Inner Detector

The ATLAS ID is the detector closest to the beam pipe and designed to perform measurements on charged particles originating directly from the interaction point. It combines high momentum resolution of  $\sigma_{p_T}/p_T = 0.05\% \oplus 1\%$  with tracking capabilities, all immersed in a solenoidal magnet field perpendicular to the beam pipe generated by the CS.

The primary purpose of the ID is to track trajectories of charged particles as they travel through this subdetector; the points in space provided by the ID allow for the precise momentum resolution. Closest to the beam pipe are two silicon based detectors; the Pixel detector and the Semiconductor Tracker. These detectors are silicon based, offer the best resolution per particle hit and are close to the interaction point. The outermost layer contains the Transistor Radiation Tracker (TRT), which carries less precise but more affordable technology in the form of gaseous straw tubes. A schematic drawing of the central ID and the complete ID, including endcap layers, are shown in Fig. 6.2 and Fig. 6.3 respectively.

### 6.2.1 The Pixel Detector

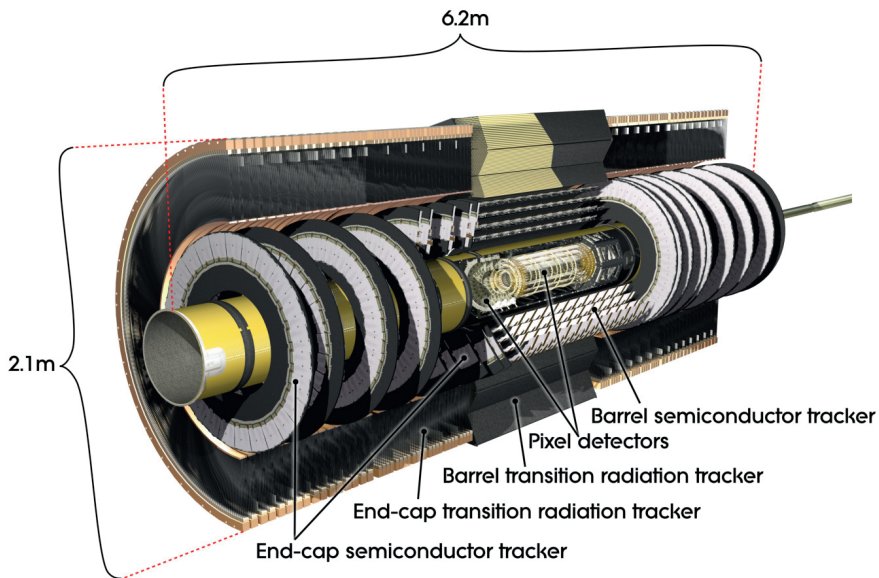
The Pixel detector [143] is illustrated in Fig. 6.2 in green, and in Fig. 6.3; it consists of four module layers in the barrel regions and three disks in each of the two end-cap structures.



**Figure 6.2:** Schematic drawing of the central subdetectors of the ATLAS ID. Showing a charged particle traversing the ID. Taken from Ref. [142].

The layers carry a total of 1744 modules, consisting of about 80 million readout channels which result in a coverage of  $|\eta| < 2.5$ . The innermost barrel layer is called the Pixel B-layer; the others are referred to as Layer-1 and Layer-2. Each pixel module is composed of 16 front-end electronics chips with 2880 channels and n-type pixel sensors of  $50 \mu\text{m} \times 400 \mu\text{m}$  with a thickness of  $250 \mu\text{m}$ .

The resolution of particle hits on the modules is  $115 \mu\text{m}$  in the longitudinal ( $\theta$ ) direction and  $10 \mu\text{m}$  transverse to the x-y plane. The pixel detector performs vertex reconstruction, which is instrumental in flavour tagging. Because one of the most abundant backgrounds within ATLAS is the  $t\bar{t}$  process, which decays into two  $b$ -quarks, flavour tagging is crucial for analyses performed in this thesis. Since this detector is closest to the interaction point, severe radiation damage to this detector is unavoidable. For this reason, the insertable b-layer acts as an extension and additional radiation shield to the pixel detector.



**Figure 6.3:** Schematic drawing of the ATLAS ID. Taken from Ref. [142].

### 6.2.2 The Insertable $b$ -layer

The Insertable  $b$ -layer [144] was added between the interaction point and the Pixel detector during the years 2013 and 2014. After it was inserted, this subdetector was mounted closest to the beam pipe at about a radial distance of 33 mm. As the name suggests, it was added because ATLAS was losing important tracking information, specifically about the  $b$ -quark, due to the radiation damage during runtime. The  $b$ -layer improves the determination of the track parameters and also the capability of the reconstruction of secondary vertices, crucial for flavour tagging algorithms. The pixel size in the insertable  $b$ -layers is  $50\text{ }\mu\text{m} \times 250\text{ }\mu\text{m}$  and provides a resolution of  $47\text{ }\mu\text{m}$  in the longitudinal direction and  $8\text{ }\mu\text{m}$  in the transverse direction. This enhances tracking performance and radiation tolerance, with the tracker covering pseudorapidities up to  $|\eta| < 3$ .

### 6.2.3 The Semiconductor Tracker

The Semiconductor Tracker [145] is a silicon strip detector located around the pixel detector. In order to provide a two-dimensional measurement in the module plane of charged particles originating from the beam-interaction region, the Semiconductor Tracker measures four space-points. This allows the ID to establish a measurement of the  $z$ -coordinate along the strip length and improves spatial resolution in conjunction with the  $r - \phi$  measurement.

Shown in Fig. 6.2 (in blue) and Fig. 6.3, it consists of four cylindrical layers in the barrel region and nine disks in the end-cap region. The subdetector is composed of 4088 modules,

each of which consists of two back-to-back micro-strip detectors with a pitch of  $80\text{ }\mu\text{m}$ . Additionally, they are tilted with respect to each other by  $40\text{ mrad}$ .

### 6.2.4 The Transistion Radiation Tracker

The TRT plays a role in particle identification. When energetic particles cross the material boundaries, which have different dielectric constants, photons are emitted; this effect is known as transition radiation. This radiation is absorbed by the TRT tubes, causing higher readout signals. Since electrons are more likely to exhibit transition radiation, because of their larger gamma factor  $\gamma = E/m$  than charged hadrons, they have a larger probability of generating TRT hits. This not only aids the electron track reconstruction, but also allows better separation of electrons and charged hadrons. On average, charged particles leave about 36 hits in the TRT, which means it substantially contributes to the overall momentum determination, despite its lower resolution.

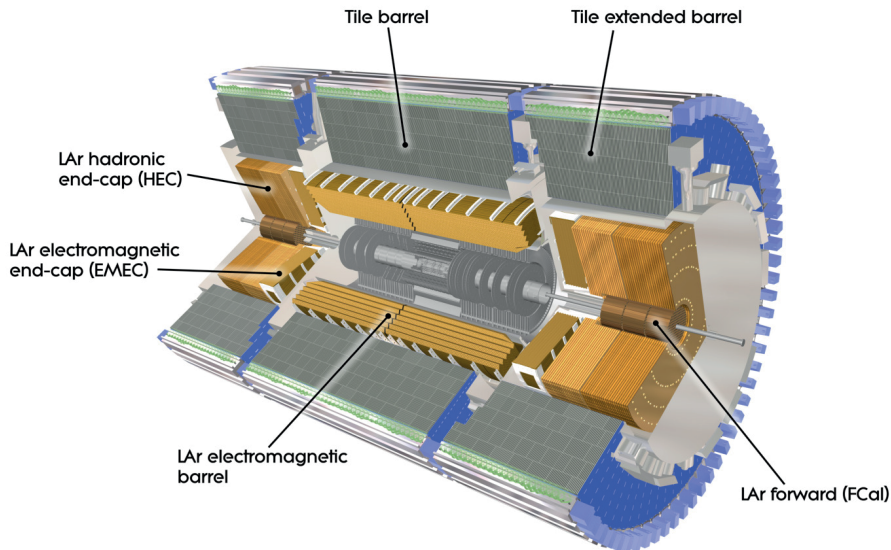
The main component of the TRT is the drift tubes [146] and is the outermost tracker of the ID, shown in Fig. 6.2 in yellow. The drift tubes are filled with a gas mixture of xenon  $Xe(70\%)$ , carbon dioxide  $CO_2(27\%)$  and oxygen  $O_2(3\%)$ . In these drift tubes, an anode is placed made of gold-plated tungsten wire. When charged particles move through these tubes, the gas gets ionized and charges are collected by the anodes, which allows for the determination of the passage of a charged particle, with a spatial resolution of  $130\text{ }\mu\text{m}$ . The drift tubes in the barrel are  $144\text{ cm}$  in length and situated parallel to the beam direction; the  $37\text{ cm}$  long drift tubes in the endcaps are present in the radial direction. The different layers are interleaved with a radiator medium, measuring the emitted transition radiation allows for mass and energy measurements of the interacting particles.

## 6.3 The calorimeter system

Once particles leave the ID, they enter the calorimeter system of the ATLAS experiment [147] depicted in Fig. 6.4, which absorb and measure the energy of particles. The complete calorimeter system consists of several subdetectors: the LAr electromagnetic calorimeter [148] (divided into the barrel and endcap parts), the Hadronic Tile Calorimeter [149], the LAr Hadronic End-Cap Calorimeter [148] and the LAr forward calorimeter [148]. These calorimeters cover the pseudorapidity range up to  $|\eta| < 4.9$  using different technologies suited to the varying physics processes over the  $\eta$  range. The fine granularity of the EM calorimeter in the  $\eta < 2.5$  range, matched to the ID, is suited for precision measurements of electrons and photons. The coarser granularity, of the rest of the calorimeters, is best suited for jet and  $E_T^{\text{miss}}$  reconstruction.

The combination of the calorimeter systems must provide good containment of electromagnetic and hadronic energy, and must also limit the punch-through of these showers towards the muon calorimeter system. Hence, the radiation thickness of this system is an important design consideration.





**Figure 6.4:** An illustration of the ATLAS calorimeter system. Taken from Ref. [147].

The total radiation thickness of the EM calorimeter is  $> 22X_0$  in the barrel region and  $> 24X_0$  in the endcaps, where  $X_0$  is one standard radiation length. The nuclear interaction length ( $\lambda$ ) of about  $10\lambda$  in both the barrel and endcaps provides good resolution for high energetic jets. The total thickness is sufficient to limit the punch-through into the muon system, which (together with the large  $\eta$  coverage) is setup to provide good measurements of  $E_T^{\text{miss}}$ . The transverse missing energy is an important signature for many physics analysis with neutrinos in the final state, which includes  $H \rightarrow WW^*$  decays.

### 6.3.1 The Electromagnetic Calorimeter

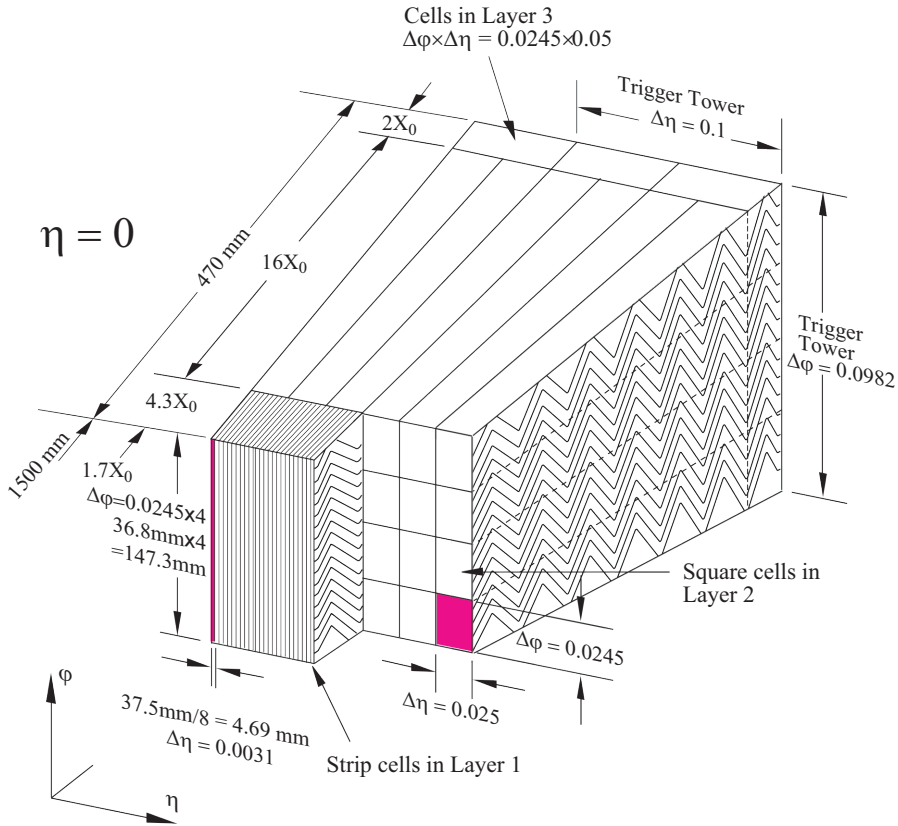
The component of the calorimeter system closest to the interaction point is the Electromagnetic Calorimeter (ECal). The ECal uses lead as an absorbing material and LAr as the active material, which converts the absorbed energy of a traversing particle into a measurable signal. The ECal is designed to measure electrons and photons. Electrons passing through the active material lose their energy by emitting photons in the bremsstrahlung process, while photons lose their energy in pair production. When absorbing electromagnetic showers, the amount of absorbed energy scales with the radiation length ( $\chi_0$ ) in the material.

The lead plates in the ECal are organized in a accordion shape, extending the radiation length substantially, which also provides full coverage in  $\phi$  without any cracks, in combination with fast extraction of the signal at the electrodes. Because of the use of heavy elements such as tungsten, copper or lead, it favours the electromagnetic



interaction of particles while keeping the nuclear interaction lengths ( $\lambda$ ) low, such that hadronic particles pass through to the Hadronic Calorimeter (HCal).

Fig. 6.5 shows the structure of an ECal section, with the three different granularities becoming less coarse when moving away from the interaction point. The first and last layers absorb the tip and tail end of the electromagnetic shower, while the middle part absorbs the largest fraction of the EM energy.



**Figure 6.5:** Sketch of a barrel section of the ECal showing the three different granularities. Taken from Ref. [147].

As shown in Fig. 6.4, the ECal is divided into three parts. The barrel part covers the  $|\eta| < 1.475$  range; the two end-cap components cover the forward region  $1.375 < |\eta| < 3.2$ . The granularity of the first layer of the calorimeter endcap is coarser than that of the central part, decreasing from  $\Delta\eta \times \Delta\phi = 0.025 \times 0.1$  to  $\Delta\eta \times \Delta\phi = 0.1 \times 0.1$ . This is also the case for the other layers of the ECal and provides enough resolution to satisfy the physics requirement for jet and  $E_T^{\text{miss}}$  reconstruction. Additionally, in the region of  $|\eta| < 1.8$ , a pre-sampler detector is placed to correct for the energy lost by electrons and photons by passing through the ID and its solenoid magnet. When these elements are combined, the ECal is able to measure energies with a resolution of

$\sigma_E/E = 9\%/\sqrt{E} \oplus 0.3\%$ . The first term in this uncertainty accounts for the stochastic uncertainty when counting photons, the second term gives the  $p_T$ -dependent uncertainty.

### 6.3.2 The Hadronic Calorimeter

After the ECAL has absorbed the energy of the electromagnetic showers, the particles travel to the HCal. The HCal measures hadronic showers by exploiting the nature of hadrons, which lose their energy by inelastic interaction with other nuclei. Similarly as for the ECAL's radiation length, in the HCal the number of nuclear interaction lengths ( $\lambda$ ) is important to infer the energy of the particle.

When hadrons lose energy through nuclear interactions, photons are measured in the active material. Therefore, the material used is dense, to ensure large amount of interactions lengths:  $\lambda$  is the mean distance a particle has to travel before a nuclear interaction.

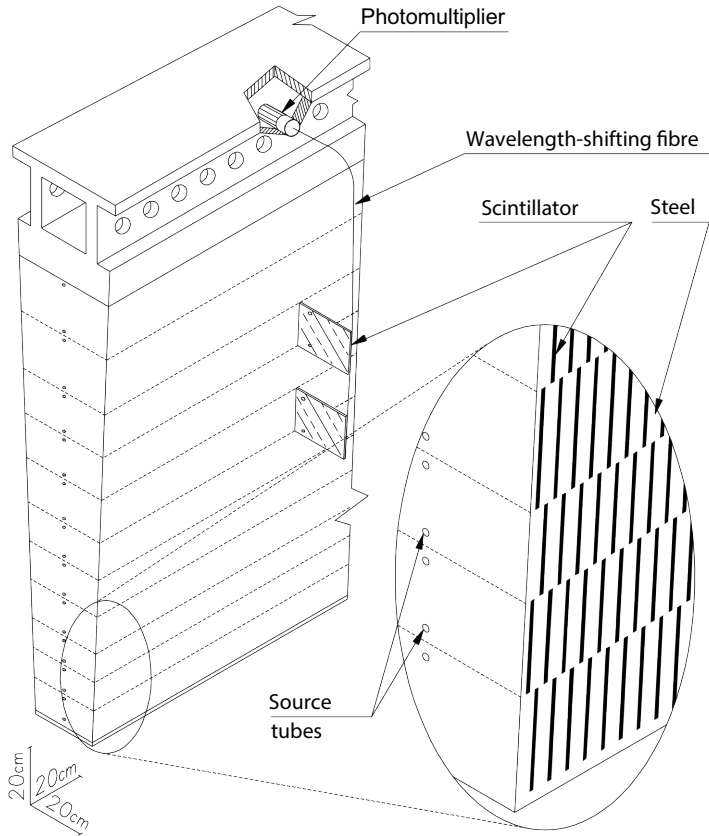
The total HCal system consists of three different subdetectors: the hadronic endcap, the Tile Calorimeter and the LAr forward calorimeter (FCal). The hadronic endcap covers  $1.5 < |\eta| < 3.2$  range; the Tile Calorimeter covers the entire barrel region  $|\eta| < 1.0$  and also extends into the endcap region to  $|\eta| < 1.7$ ; the LAr FCal covers the  $3.1 < |\eta| < 4.9$  range. Like the ECAL, the hadronic endcap also uses LAr as an active material, but uses copper as absorber material.

### 6.3.3 The Tile Calorimeter

The Tile Calorimeter, schematically drawn in Fig. 6.6, is placed directly outside the EM calorimeter envelope, providing  $9.7 \lambda$  detector thickness at  $\eta = 0$ . The self-supporting barrel and extended barrels are divided into 64 azimuthal modules. The Tile Calorimeter consists of plastic scintillating tiles as active material and uses steel as absorber with an interaction length of  $7.5 \lambda$ , providing near perfect coverage in  $\phi$ . The energy resolution of the Tile Calorimeter and the Hadronic Endcap Calorimeters combined is  $\sigma_E/E = 50\%/\sqrt{E} \oplus 0.3\%$ , using a total of 10000 readout channels.

### 6.3.4 The Forward Calorimeter

The forward calorimeter is completely integrated into the end-cap cryostats. It provides about  $10 \lambda$  interactions lengths and consists of three modules. The first absorbing material is made of copper and is optimized for EM measurements. The remaining two absorbing materials are made of tungsten and are used to predominantly measure the energy of hadronic interactions in the forward region. Each module in the forward calorimeter consists of a metal matrix with regularly spaced longitudinal channels, filled with the electrode structure of concentric rods and tubes parallel to the beam axis. The LAr in the gaps between the rod and tubes is the active medium. The energy resolution for the FCal is  $\sigma_E/E = 100\%/\sqrt{E} \oplus 10\%$ , using 1800 readout channels.

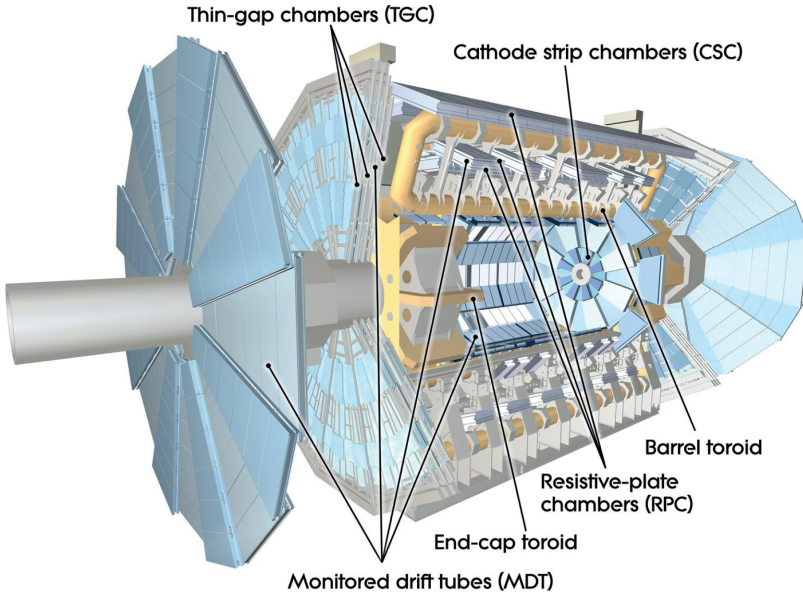


**Figure 6.6:** Schematic drawing of the assembly of the Tile Calorimeter, including the optical readout components. Taken from Ref. [147].

## 6.4 The Muon Spectrometer

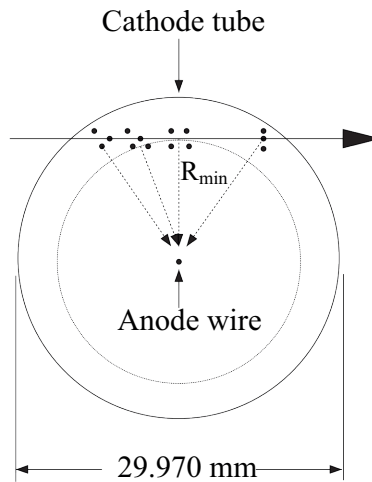
The MS [150], illustrated in Fig. 6.7, forms the outermost part of the ATLAS detector and is designed to detect muons or muon-like charged particles exiting the barrel and the end-cap calorimeter system. Using the magnetic field generated by the ATLAS Toroid system, it measures the momentum of the charged particles from the bent tracks in the pseudorapidity range of  $|\eta| < 2.7$ . The muon system's goal is the stand-alone measurement of the transverse momentum and to trigger on highly energetic tracks.

The precision momentum measurement is performed by the Monitored Drift Tube (MDT) chambers. These tubes provide a measurement of the track coordinates over almost the entire  $\eta$  range, designed for high accuracy, simplicity, and predictability in mechanical deformations. The MDTs cover the pseudorapidity range  $|\eta| < 2.7$ , except for the innermost endcap layer, where the coverage is  $|\eta| < 2.0$ . The chambers consist of three to eight layers of drift tubes and achieve a resolution of  $80 \mu\text{m}$  per tube, or  $35 \mu\text{m}$  per chamber when assembled in groups.



**Figure 6.7:** Schematic drawing of the ATLAS muon subsystem [147].

An MDT, shown in Fig. 6.8, is a 30 mm diameter aluminium drift tube with a maximum drift time of 700 ns. When a charged particle passes through the Argon (Ar)/Carbon dioxide ( $CO_2$ ) gas operating at 3 bar, the gas becomes ionized; the resulting electrons from such ionization are collected at the central tungsten-rhenium anode wire by an electric potential. The position of the wire inside the tube, placed with an accuracy of  $\sigma < 10 \mu m$ , is used to infer the coordinate of the passed charged particle.

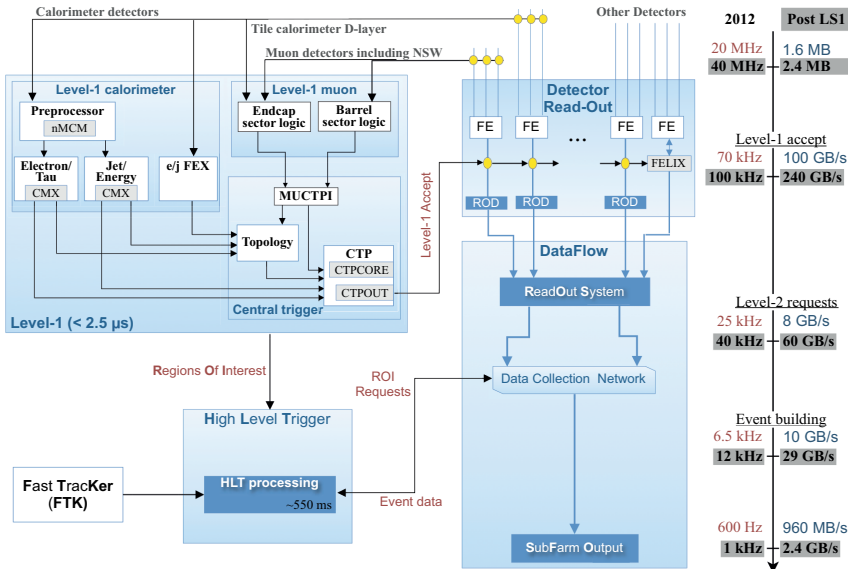


**Figure 6.8:** Cross-section of an MDT tube with a charged particle passing through, creating electrons via ionization. Taken from Ref. [147].

At larger pseudorapidity,  $2 < |\eta| < 2.7$ , Cathode Strip Chambers (CSCs) are used, with a higher granularity in order to accommodate for the higher background rate. The CSCs are trapezium-shaped wire chambers and consist of an anode wire and two cathode strips, which provide a resolution of  $40\ \mu\text{m}$  in the  $\eta$ -plane, the bend-plane, and  $5\ \text{mm}$  in the  $\phi$ -plane, the non-bending-plane. A low maximum electron drift time of  $40\ \text{ns}$  allows for a timing resolution of  $7\ \text{ns}$  per plane. Combined the MDTs and CSCs allow for high precision measurement of the momentum and trajectory: the momentum resolution uncertainty is below  $3\%$  with a  $p_T$  between  $10$  and  $100\ \text{GeV}$  increasing to about  $10\%$  for muons with a  $p_T$  of  $1\ \text{TeV}$ .

## 6.5 Triggering and data acquisition

The Trigger system in Run-2 [151, 152], shown in Fig. 6.9, is designed to only select the events passing the selections defined by the triggers within the ATLAS Trigger menu [153]. This predefined set of criteria decides if an event is kept for further storage and processing based on a chain of trigger selections. The bunch spacing of  $25\ \text{ns}$  translates into a collision rate of  $40\ \text{MHz}$ , resulting in a challenging environment to separate the rare and interesting events from the large amount of background processes. Since the online event storage is limited to approximately  $1\ \text{kHz}$ , corresponding to a writing rate of  $2.4\ \text{GB s}^{-1}$ , the events are selected and the rate is reduced by a factor of  $10^4$  through the continuously evolving Trigger and Data Acquisition System (TDAQ).



**Figure 6.9:** Schematic overview of the Phase-I Trigger and DAQ system as used in Run-2. Taken from Ref. [152].

The ATLAS trigger system can be viewed as a two stage system consisting of the Level 1 (L1) trigger and the High Level Trigger (HLT). At L1, the bulk of the reduction takes

place: the 40 MHz rate is brought down to roughly 100 kHz with a latency of 2.5 ns. By a reprogrammable hardware trigger system; the algorithms run at low-latency on Field Programmable Gate Arrays (FPGAs). The second level, the HLT, is a software-based trigger where the algorithms run at computing servers. The HLT algorithms have access to the full event information and further reduce the rate to the required 1 kHz, with a latency of the order of seconds. The complete TDAQ system is responsible for the transport, rate reduction, assembly and storage of the event data from the subdetectors.

### 6.5.1 L1 Trigger

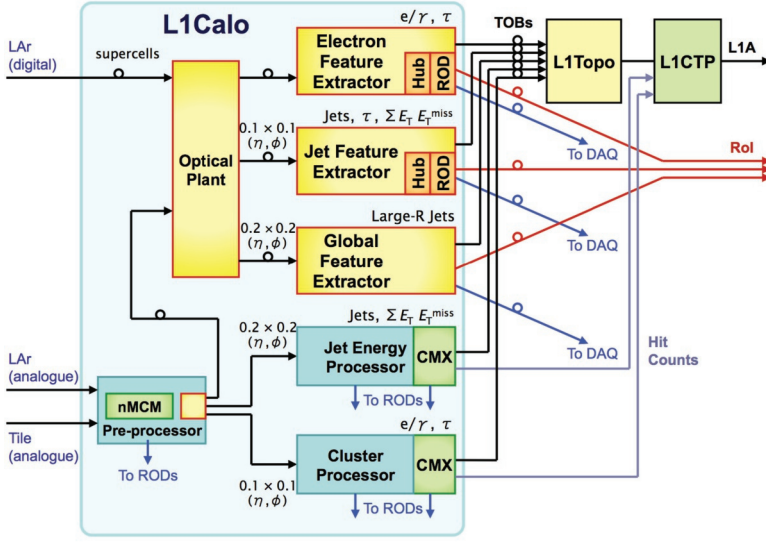
In Run-2, the L1 Trigger has three subsystems: The L1Calo, L1Muon and L1Topo. In this setup, the L1Calo and L1Muon systems only have access to the calorimeter and muon event information, respectively. L1Topo is a newer addition to the system: it has access to both systems and is able to run more complicated algorithms using several types of Trigger Objects (TOBs). The L1 system consists of built-in electronics that compares incoming signals to preset programmable items. If any threshold for at least one item is met, a L1 accept action is issued; this decision is made by the Central Trigger Processor (CTP). The L1 trigger that fires the readout starts the processing of the event data in the Readout Drives (ROD).

#### L1Calo

The legacy L1Calo system, used during Run-2 and depicted in green in Fig. 6.10, receives input from roughly 7200 analogue trigger towers of the electromagnetic and hadronic calorimeters, with a granularity of  $\Delta\eta \times \Delta\phi = 0.1 \times 0.1$ . The first Module is the nMCM, which digitizes, calibrates and filters the calorimeter signals. This module also serves as a Preprocessor (PP) which assigns a Bunch Crossing ID (BCID) to the signal. Furthermore, pedestal corrections, noise suppressions and energy calibration are applied, before transferring the data to the Cluster Processor (CP) and the Jet/Energy processor (JEP). The JEP executes the jet clustering, while the CP is tasked with finding electrons, photons and  $\tau$  candidates. Both processors transmit the information about passed thresholds to the CTP. Three new feature extractors (FEX): the electromagnetic (eFEX), jet (jFEX) and global (gFEX) feature extractors, are added to the L1Calo system in the Phase-I upgrades [154]. These modules, depicted in yellow in Fig. 6.10, gradually take over the functions of the legacy processing sub-systems during Run-3.

The three new FEX systems process different input information and feature different algorithmic trigger requirements, which are separately optimized for each of the trigger signatures. The eFEX is responsible for the electron and photon trigger signatures; the jFEX for the jet trigger signatures and  $E_T^{\text{miss}}$ . The gFEX module is specialized on global trigger signatures, along with large-R jets.

In addition to the new modules, the granularity of the trigger towers are also improved and digitized by the replacement with Super Cells. The finer granularity is designed



**Figure 6.10:** Schematic overview of the L1Calo system after the Phase-I Upgrades of the ATLAS TDAQ System which includes the legacy L1Calo system used for Run-2 in addition to the Run-3 gFex and jFex feature extractor modules [154].

to improve the performance of the Level-1 single object triggers by enabling the use of shower-shape variables for more effective identification of electrons, photons, and  $\tau$  leptons, sharpening the EM, jet, and  $E_T^{\text{miss}}$  efficiency turn-on curves using more advanced reconstruction algorithms at Level-1 and utilizing pileup subtraction techniques on an event-by-event basis.

The old trigger tower granularity is compared to the Phase-I upgraded super cell granularity in Tab. 6.1. Several super cells are merged into trigger towers with varying granularity, referred to as jTowers and gTowers for the jFEX and gFEX modules, respectively; these are used to compute inputs for trigger decisions.

**Table 6.1:** Comparison of the Run-2 Trigger Tower granularity vs. the Phase-I upgraded Super Cell granularity in the LAr EM barrel calorimeter. The number of elementary cells grouped for the trigger readout in  $\eta$  and  $\phi$  are indicated by  $n_\eta$  and  $n_\phi$  [154].

| Layer | Elementary Cell | Trigger Tower          |                                | Super Cell             |                                |
|-------|-----------------|------------------------|--------------------------------|------------------------|--------------------------------|
|       |                 | $n_\eta \times n_\phi$ | $\Delta\eta \times \Delta\phi$ | $n_\eta \times n_\phi$ | $\Delta\eta \times \Delta\phi$ |
| 0     | Presampler      | $4 \times 1$           | $0.025 \times 0.1$             | $4 \times 1$           | $0.1 \times 0.1$               |
| 1     | Front           | $32 \times 1$          | $0.003125 \times 0.1$          | $8 \times 1$           | $0.025 \times 0.1$             |
| 2     | Middle          | $4 \times 4$           | $0.025 \times 0.025$           | $1 \times 4$           | $0.025 \times 0.1$             |
| 3     | Back            | $2 \times 4$           | $0.05 \times 0.025$            | $2 \times 4$           | $0.1 \times 0.1$               |



## L1Muon

The L1Muon system is completely separate from the L1Calo system and processes the information originating from the MS and the muon trigger chambers: the TGCs and RPCs. In total, there are over 800000 input channels to the L1Muon system; the timing resolution of the MS is good enough to allocate the muon candidates to a BCID. A high- $p_T$  muon candidate is required to pass all three RPC stations in the barrel and all TGCs in the forward region; for low- $p_T$  muons only, signal deposition in the first two stations is required.

## L1Topo

The L1 topological trigger, commissioned during 2015/2016, enables the combination of TOBs from the L1Calo and L1Muon systems. It achieves an energy resolution of 1 GeV and is the first trigger module in the chain to incorporate the TOBs'  $\eta$  and  $\phi$  information. This capability permits a more complex selection and combination of TOBs via more sophisticated trigger algorithms compared to L1Calo and L1Muon systems alone. In particular, the L1Topo system bases its decisions on the two leading muon candidates, as well as the electron, photon,  $\tau$ , and jet candidates, in addition to the  $E_T^{\text{miss}}$ .

### 6.5.2 HLT

The final trigger algorithms the event sees before it is saved are the HLT algorithms. The HLTs operate solely on the L1-passed TOBs, which serve as the L1 seed. They are designed to closely mirror the offline event reconstruction to ensure reliable trigger efficiency. Once an HLT event is seeded, the event building step ensues, and the trigger algorithm reconstructs the event based on predefined requirements. At last, when the sequence of L1 seeding and HLT verification is complete, the event successfully passes through the entire trigger chain.

### 6.5.3 Triggers used in $H \rightarrow WW$ analyses

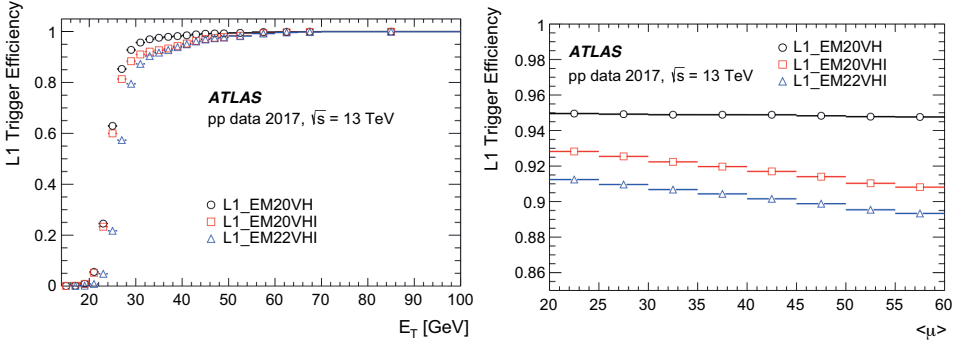
The  $H \rightarrow WW$  analyses in this thesis select events via a combination of different-flavour di-lepton and single-lepton triggers. The di-lepton triggers require one electron and one muon, each exceeding a specified  $p_T$  threshold. In contrast, single-lepton triggers require either an electron or muon with a higher  $p_T$  threshold to ensure an acceptable trigger rate.

The triggers listed in [Tab. 6.2](#) are named according to their  $E_T$  and  $p_T$  thresholds, as well as their identification and isolation requirements. Triggers with lower  $p_T$  thresholds generally have stricter identification and isolation criteria. They are also categorized by data-taking periods, reflecting changes in the triggers used over time.

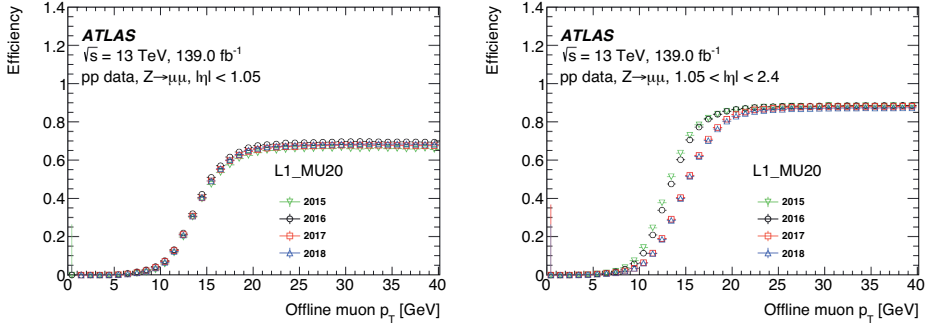
The on-shell  $H \rightarrow WW$  measurement uses a trigger combination that maximizes signal efficiency by combining base single-lepton triggers with di-lepton triggers to improve efficiency in the low- $p_T$  region. The trigger efficiency for two key L1 triggers is shown in



Fig. 6.11 for electron triggers and Fig. 6.12 for muon triggers.



**Figure 6.11:** Performance of the L1 Electron triggers during Run-2. Left: the efficiency as a function of  $E_T$ . Right: the efficiency as a function of the pile-up [155].



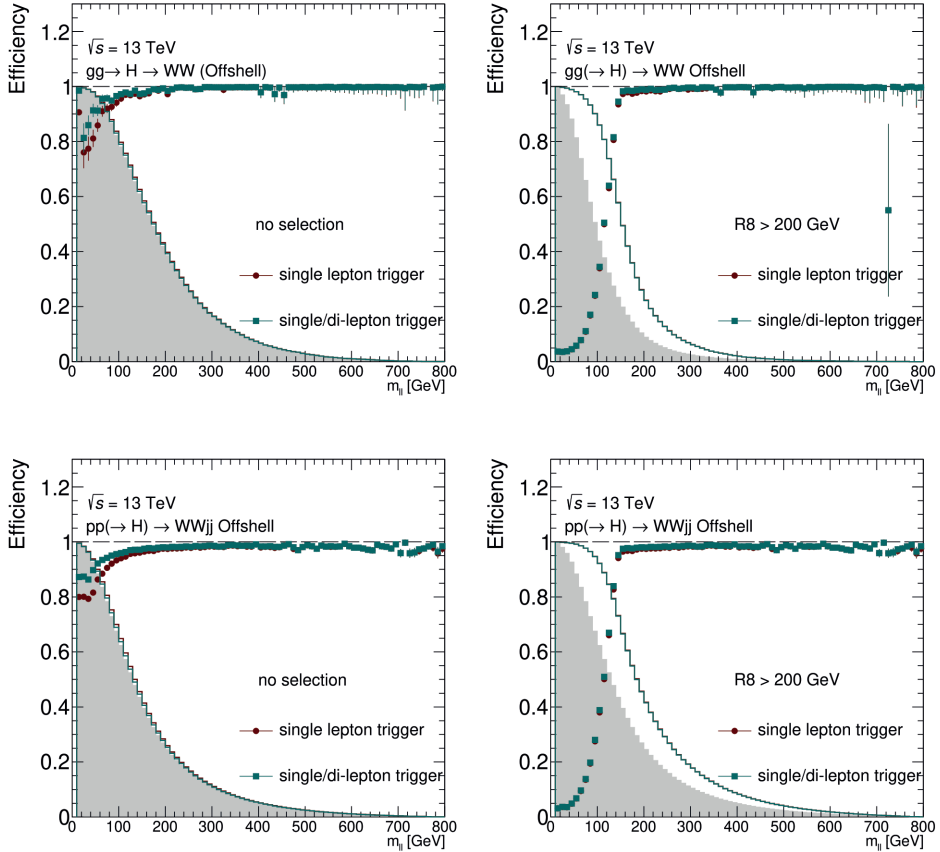
**Figure 6.12:** Efficiency of the L1 Muon triggers during Run-2. Left: the efficiency in the barrel. Right: the efficiency in the endcaps [156].

Unlike the on-shell analysis, the off-shell  $H^* \rightarrow WW$  analysis uses only single-lepton triggers, as di-lepton triggers do not improve efficiency, as shown in Fig. 6.13. This is because the condition  $\sqrt{m_{ll}^2 + (0.8 * m_T)^2} > 200 \text{ GeV}$ <sup>1</sup> is typically satisfied and excludes additional low- $p_T$  leptons accepted by di-lepton triggers.

To include effects from the usage of trigger, Monte Carlo samples are corrected with scale factors in order to account for the differences in efficiency between simulated and real data.

Additionally, simulated samples are corrected for the potential pre-scale of triggers, which differs between data-taking periods; the pre-scale-factor is in the range of 1-8798 for both single electron and muon triggers used in the estimate of mis-identified leptons, labelled as the  $Z$ +jets and di-jets fakes in Tab. 6.2. Only unprescaled triggers are used for the signal selection in this thesis.

<sup>1</sup> The transverse mass  $m_T$  is defined as  $m_T = \sqrt{(E_{ll} + E_T^{\text{miss}})^2 - |p_{ll} + E_T^{\text{miss}}|^2}$



**Figure 6.13:** Comparison of signal efficiency of di-lepton/single-lepton triggers vs. single-lepton triggers as a function of the transverse mass  $m_T$ . Before (left) and after (right) applying  $\sqrt{m_{ll}^2 + (0.8 * m_T)^2} > 200$  GeV. The signal efficiency is shown for the ggF signal topology (top) and the VBF signal topology (bottom).

**Table 6.2:** Trigger chains used in the  $H \rightarrow WW$  di-jets fakes,  $Z$ +jets fakes, on-shell and off-shell analyses. Multiple L1 triggers indicate different seeds for specific data-taking periods. Numbers after ‘EM’, ‘MU’, ‘e’ and ‘mu’ denote  $E_T$  and  $p_T$  thresholds at L1 and HLT levels, respectively. The terms *tight*, *medium*, *loose* describe ID criteria; *i* denotes isolation.

| Lepton    | Level-1 Trigger  | High Level Trigger             | Used in $H \rightarrow WW$ Analysis  |
|-----------|------------------|--------------------------------|--------------------------------------|
| 2015      |                  |                                |                                      |
| $e$       | L1_EM10VH        | HLT_e12_lhloose_L1EM10VH       | di-jets fakes                        |
|           | L1_EM20VH        | HLT_e24_lhmedium_L1EM20VH      | on-shell, off-shell                  |
|           | L1_EM22VHI       | HLT_e60_lhmedium               | $Z$ +jets fakes, on-shell, off-shell |
|           | L1_EM22VHI       | HLT_e120_lhloose               | $Z$ +jets fakes, on-shell, off-shell |
| $\mu$     | L1_MU10          | HLT_mu14                       | di-jets fakes                        |
|           | L1_MU10          | HLT_mu24                       | di-jets fakes                        |
|           | L1_MU10          | HLT_mu50                       | di-jets fakes                        |
|           | L1_MU15          | HLT_mu20_iloose_L1MU15         | $Z$ +jets fakes, on-shell, off-shell |
|           | L1_MU20          | HLT_mu50                       | $Z$ +jets fakes, on-shell, off-shell |
| $e\mu$    | L1_EM15VH_MU10   | HLT_e17_lhloose_mu14           | on-shell                             |
| 2016–2018 |                  |                                |                                      |
| $e$       | L1_EM10VH        | HLT_e12_lhloose_nod0           | di-jets fakes                        |
|           | L1_EM22VHI       | HLT_e26_lhtight_nod0_ivarloose | on-shell, off-shell                  |
|           | L1_EM22VHI       | HLT_e60_lhmedium_nod0          | $Z$ +jets fakes, on-shell, off-shell |
|           | L1_EM22VHI       | HLT_e140_lhloose_nod0          | $Z$ +jets fakes, on-shell, off-shell |
| $\mu$     | L1_MU10          | HLT_mu14                       | di-jets fakes                        |
|           | L1_MU10          | HLT_mu24                       | di-jets fakes                        |
|           | L1_MU10          | HLT_mu50                       | di-jets fakes                        |
|           | L1_MU20, L1_MU21 | HLT_mu26_ivarmedium            | $Z$ +jets fakes, on-shell, off-shell |
|           | L1_MU20, L1_MU21 | HLT_mu50                       | $Z$ +jets fakes, on-shell, off-shell |
| $e\mu$    | L1_EM15VH_MU10   | HLT_e17_lhloose_nod0_mu14      | on-shell                             |

## RUN-3 UPGRADES FOR L1 MISSING TRANSVERSE MOMENTUM TRIGGERS

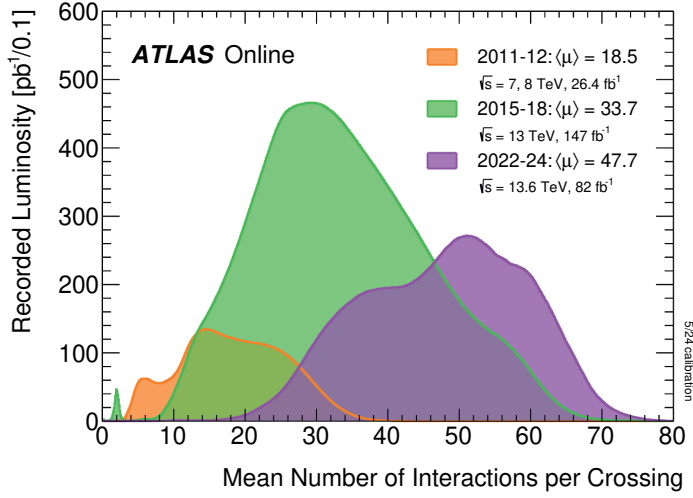
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This chapter describes methods that apply online L1 jet-based corrections, determined using offline reconstructed jet and  $E_T^{\text{miss}}$  information, to the L1 reconstructed  $E_T^{\text{miss}}$ . These methods aim to enhance the performance of low-level  $E_T^{\text{miss}}$  algorithms under high pile-up conditions and maintain a stable trigger rate in Run-3 and beyond. Additionally, the performance of a new Run-3 L1  $E_T^{\text{miss}}$  algorithm deployed on the gFEX module is optimized using a Kalman filter. Finally,  $E_T^{\text{miss}}$  trigger algorithm efficiencies and trigger rates are estimated using Run-3 emulated data and Monte Carlo simulations.

Invisible particles, that neither interact via the strong or electromagnetic force nor decay inside the detector, leave no signature. Yet, triggering on events with these particles is crucial for many ATLAS studies. For instance, the SM Higgs decay into  $b$ -quarks (included in the SM-EFT interpretation in [Chapter 12](#)) was initially observed in ZH production, where the Z boson decays into neutrinos. Other examples are searches for invisible Higgs decays [\[157\]](#) and supersymmetric scenarios with an invisible neutralino [\[158\]](#).

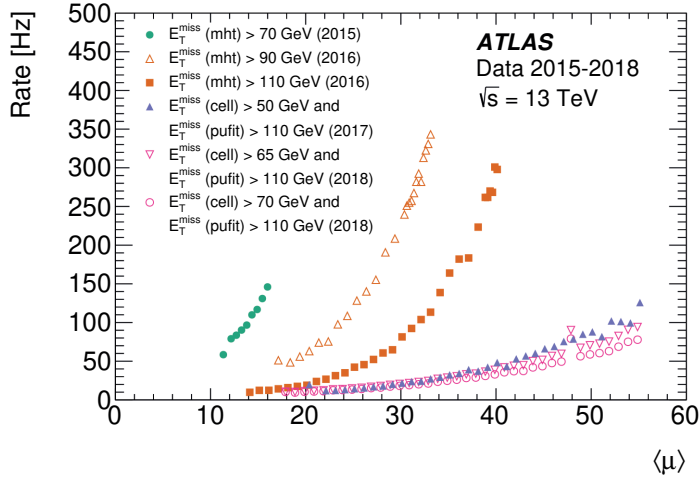
In addition to enhancing trigger efficiency for invisible particles, the rising centre-of-mass energy and luminosity in each LHC run necessitate further developments in missing transverse momentum triggers to maintain the required performance. LHC Run-1 operated from 2010 to 2012 at a centre-of-mass energy of 7 and 8 TeV. The second run operated from 2015 to 2018 at  $\sqrt{s} = 13$  TeV with peak luminosities reaching  $1.9 \times 10^{34} \text{ cm}^{-2} \text{ s}^{-1}$  [\[137\]](#). With the LHC progressing through Run-3, it is aiming for even higher luminosities; increasing the number of  $pp$ -collisions within each bunch crossing and results in bigger challenges.

The additional collisions facilitate the increased luminosity of the LHC; this increase in the mean number of interactions over the years, indicated as mean pile-up  $\langle\mu\rangle$ , is shown in [Fig. 7.1](#). During Run-2, a peak pile-up of approximately 70 interactions per bunch crossing was observed. The energy from these additional pile-up collisions is deposited throughout the detector, degrading the  $E_T^{\text{miss}}$  resolution because of the extra hadronic activity in the bunch crossing. Failure to account for this pile-up can lead to an unacceptable order-of-magnitude increase in trigger rates, unless  $E_T^{\text{miss}}$  trigger thresholds are raised, which would consequently decrease signal efficiency.



**Figure 7.1:** The luminosity-weighted distribution of the mean number of interactions per crossing for pp collisions at 7 and 8 TeV centre-of-mass energy in 2011-12, 13 TeV in 2015-18, and 13.6 TeV in 2022-24 [159].

This raise of trigger thresholds is illustrated in Fig. 7.2, where each Run-2 year, due to increases in pile-up, higher trigger thresholds are required to maintain a manageable trigger rate.



**Figure 7.2:** High-level trigger rates, as a function of the pile-up  $\langle\mu\rangle$ , shown separately for example runs in each year for 2015-2018. The triggers shown are named HLT\_xe70\_mht (2015), HLT\_xe90\_mht and HLT\_xe110\_mht (2016), HLT\_xe110\_pufit (2017), HLT\_xe110\_pufit\_xe65 and HLT\_xe110\_pufit\_xe70 (2018). The HLT\_xe110\_pufit trigger used during 2017 required cell  $E_T^{\text{miss}} > 50 \text{ GeV}$ ; all triggers are seeded by L1XE50, which requires  $E_T^{\text{miss}} > 50 \text{ GeV}$  at L1 [160].

In 2015 and 2016, a steep rise in trigger rate of the mht trigger algorithm is observed, which is a trigger that builds  $E_T^{\text{miss}}$  out of high- $p_T$  jet activity. This steep increase was remedied by tightening the trigger thresholds with the introduction of the cell algorithm, which builds  $E_T^{\text{miss}}$  out of the calorimeter cell information, and by introducing the pufit algorithm in 2017 [160].

The pufit algorithm corrects for pile-up effects on high- $E_T$  calorimeter signals contributing to  $E_T^{\text{miss}}$ . It produces a fit, based on topological patches of the detector in  $\eta$  and  $\phi$ , which is used to estimate the energy originating from pile-up. After pile-up subtraction, the pile-up free energy patches are used to estimate the  $E_T^{\text{miss}}$  in the event. This pufit algorithm was used in 2017 and 2018 to maintain a stable trigger rate.

## 7.1 L1 missing transverse momentum trigger algorithms

The L1 Trigger system is part of the ATLAS TDAQ and serves as the first layer of event selection in the ATLAS experiment explained in § 6.5.1. The L1 triggers are tailored to prioritize low-latency, rapid processing, and high background rejection in order to control the trigger rate.

The calorimeter energy used in the decision making is measured in the calorimeter cells and is linked to energy deposits called elements, indexed by  $i$  with energy  $E_i$ , polar angle  $\theta_i$ , and azimuthal angle  $\phi_i$ . The missing transverse momentum two-vector  $\vec{E}_T^{\text{miss}}$  components constructed from these elements are defined as

$$E_x^{\text{miss}} = - \sum_{i=1}^{N_{\text{elements}}} E_i \sin \theta_i \cos \phi_i$$

$$E_y^{\text{miss}} = - \sum_{i=1}^{N_{\text{elements}}} E_i \sin \theta_i \sin \phi_i.$$

The magnitude  $E_T^{\text{miss}} = \sqrt{(E_x^{\text{miss}})^2 + (E_y^{\text{miss}})^2}$  is typically used for event selection and the total transverse energy is defined using the elements as  $\sum E_T = \sum_{i=1}^{N_{\text{elements}}} E_i \sin \theta_i$ . L1 algorithms differ in element selection, energy summation, and energy correction [161]. The L1  $E_T^{\text{miss}}$  trigger algorithms considered in this chapter are:

- The *noise-cut* (XENOISECUT) algorithm, which builds  $E_T^{\text{miss}}$  from the vector sum of all elements with  $E_T$  above an  $\eta$ -dependent noise threshold.
- The  $\rho$ -cut (XERHO) algorithm, which computes  $E_T^{\text{miss}}$  from the vector sum of all elements with  $E_T$  above a threshold depending on  $\eta$  and the local per-event pileup density ( $\rho$ ).
- The *Kalman Filter*  $E_T^{\text{miss}}$  (KFMET), this algorithm applies a jet-based correction determined with offline information to the  $\rho$ -cut algorithm to enhance the performance under high pile-up conditions [162].

- The *jets-without-jets* (jwoj) algorithm computes  $E_T^{\text{miss}}$  based on a linear combination of the soft and hard contributions to the  $E_T$  of all elements [163].

As the work in this thesis was limited to the optimization and implementation of the KFMET and jwoj algorithms during Run-3, their design and optimization are explained in the next sections. Additional information on the XENOISECUT and XERHO algorithms, as well as the gFEX and jFEX modules are documented in Ref. [164].

### 7.1.1 Kalman Filters for L1 $E_T^{\text{miss}}$ triggers

The L1Topo FPGA provides the option to build a jet corrected  $E_T^{\text{miss}}$ , achieved by modifying the L1  $E_T^{\text{miss}}$  value it receives from the L1Calo system with a Kalman filter. This uses a set of weights, as a function of jet  $p_T$  and  $|\eta|$ , hardcoded into the configuration settings of the L1Topo FPGA. The purpose of this correction is to have access to a pile-up corrected  $E_T^{\text{miss}}$  at L1, which is as close as possible to the hard-scatter  $E_T^{\text{miss}}$  of the event used in offline analyses. This leads to an improved  $E_T^{\text{miss}}$  resolution and a higher trigger efficiency.

This correction is applied using the Kalman filter based algorithm implemented on the L1Topo module, which calculates a correction value from a weighted sum of light jet TOBs reconstructed at L1. The Kalman filter is an algorithm which derives the best average value of a set of weights for a series of events by taking into account the previous values of the variable and the uncertainty of the measurement. These weights are optimized for resolution and trigger efficiency using the Kalman filter and trained on Monte Carlo samples that carry significant pile-up. These weights are then stored into 7-bit precision look-up tables (LUTs) depending on the L1 jet  $p_T$  and  $|\eta|$ .

The corrected  $E_T^{\text{miss}}$  is given by

$$\vec{E}_{T,KF}^{\text{miss}} = \vec{E}_{T,L1}^{\text{miss}} - \sum_i w \times \vec{p}_{T,i}^{\text{jet}},$$

and the quantity  $Q^2$  is defined as

$$Q^2 = \left( \vec{y} - \vec{h}_0 - \mathcal{H}\vec{w} \right)^T \mathbf{V}^{-1} \left( \vec{y} - \vec{h}_0 - \mathcal{H}\vec{w} \right).$$

In these equations, the weight vector is determined by minimizing  $Q^2$  using the Kalman filter algorithm. In the definition of  $Q^2$ , the vector  $\vec{y}$  is filled with the hard-scatter offline  $E_T^{\text{miss}}$  values  $E_{T,\text{offline}}^{\text{miss}}$  using the *Tight* working point; the  $\vec{h}_0$  is given by uncorrected  $E_T^{\text{miss}}$  values  $E_{T,L1}^{\text{miss}}$ .  $\mathcal{H}$  is the matrix of jet TOBs in the LUT binning in jet  $p_T$  of 6 bins; the bins start at 3.2 GeV and end at 1638.4 GeV, which is the maximum L1Topo jet  $E_T$  due the 13 bits used to encode energy for these objects. The  $|\eta|$  binning consists of 17 bins, starting at 0 and ending at 4.9. The vector  $\vec{w}$  is the correction vector corresponding to the correction weights  $w_i$  in the LUT. The matrix  $\mathbf{V}$  is the covariance matrix of the jet TOBs and encodes the expected uncertainty.

The correction weights are determined by minimizing the quantity  $Q^2$ , which in the case of the KFMET algorithm, is defined as

$$Q^2 = \sum_{i=1}^N \left[ \frac{E_x^{\text{offline}}(i) - E_x^{L1}(i) + \sum_{j=1}^M \sum_{k=1}^{\text{jets}(i)} w_j \cdot E_x(k) \cdot \theta(E_T^{j,k}) \cdot \theta(\eta^{j,k})}{\sigma_i} \right]^2 + \sum_{i=1}^N \left[ \frac{E_y^{\text{offline}}(i) - E_y^{L1}(i) + \sum_{j=1}^M \sum_{k=1}^{\text{jets}(i)} w_j \cdot E_y(k) \cdot \theta(E_T^{j,k}) \cdot \theta(\eta^{j,k})}{\sigma_i} \right]^2.$$

In the definition of  $Q^2$ , the variable  $\theta = 1$  if  $j = k$  and  $\theta = 0$  otherwise, to ensure the correction weight  $w_j \cdot E_x(k)$  is only applied to the offline  $E_T^{\text{miss}}$  if the jet  $k$  is within the same LUT binning as the weight  $w_j$ . The iterative algorithm is written in terms of the Kalman gain  $\mathbf{k}_k$ , the weight vector  $\hat{\mathbf{w}}_k$ , and the covariance matrix  $\hat{\mathbf{C}}_k$ , with  $k$  defined to be the iteration step as

$$\begin{aligned} \mathbf{k}_k &= \hat{\mathbf{C}}_{k-1} \mathbf{h}_k \left( \sigma_k^2 + \mathbf{h}_k^T \hat{\mathbf{C}}_{k-1} \mathbf{h}_k \right)^{-1} \\ \hat{\mathbf{w}}_k &= \hat{\mathbf{w}}_{k-1} + \mathbf{k}_k \left( E_{x/y}^{\text{offline}}(k) - E_{x/y}^{L1}(k) - \mathbf{h}_k^T \hat{\mathbf{w}}_{k-1} \right) \\ \hat{\mathbf{C}}_k &= (1 - \mathbf{k}_k \mathbf{h}_k^T) \hat{\mathbf{C}}_{k-1} \end{aligned}$$

and solved iteratively. In this equation, the value of  $\sigma_k^2$  is tuned during training. The minimization is initialized using a weight vector initialized with random values between 0 and 1; once minimized, the correction weights are used to obtain the corrected  $E_{T,KF}^{\text{miss}}$  values for the L1  $E_T^{\text{miss}}$  trigger using

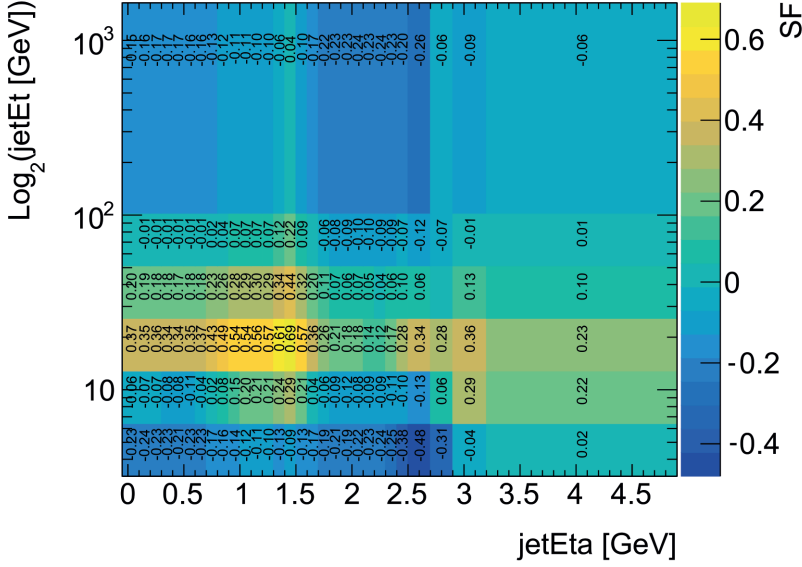
$$\begin{aligned} E_{T,KF}^{\text{miss}} &= E_T^{L1} - \sum_{i=1}^{\text{jets}} a(E_T(i), \eta(i)) \cdot E_T(i) \\ \begin{pmatrix} E_x^{\text{KF}} \\ E_y^{\text{KF}} \end{pmatrix} &= \begin{pmatrix} E_x^{L1} \\ E_y^{L1} \end{pmatrix} - \sum_{i=1}^{\text{jets}} a(E_T(i), \eta(i)) \cdot \begin{pmatrix} E_x(i) \\ E_y(i) \end{pmatrix} \\ E_{T,KF}^{\text{miss}} &= \sqrt{(E_x^{\text{KF}})^2 + (E_y^{\text{KF}})^2}. \end{aligned}$$

The best performing LUT, chosen through a process of exhausting all possible parameters and Monte Carlo combinations, is shown in [Fig. 7.3](#).

This LUT was obtained by training on a  $ZH \rightarrow \nu\bar{\nu}b\bar{b}$  sample, which contains real  $E_T^{\text{miss}}$  along with two  $b$ -jets and a  $40 \leq \mu \leq 80$ . The parameter  $\sigma_i$  was set to  $E_{T,\text{offline}}^{\text{miss}}/250$ , which enhances the relative emphasis (weight) on events with large  $E_T^{\text{miss}}$ .

Generally, it is found that the best performing LUT is obtained when the training is performed in a phase-space with a poor  $E_T^{\text{miss}}$  resolution, due to large jet activity. In the cases of large jet activity, the Kalman filter learns to correct the most poorly reconstructed events, which are most likely to be misclassified by the L1  $E_T^{\text{miss}}$  trigger, but do have real  $E_T^{\text{miss}}$ .





**Figure 7.3:** The best performing LUT for the KFMET  $E_T^{\text{miss}}$  algorithm. The minimized weights, and the scale factor (SF) they represent, are shown in each of the jet  $p_T$  and  $|\eta|$  bins they correspond to.

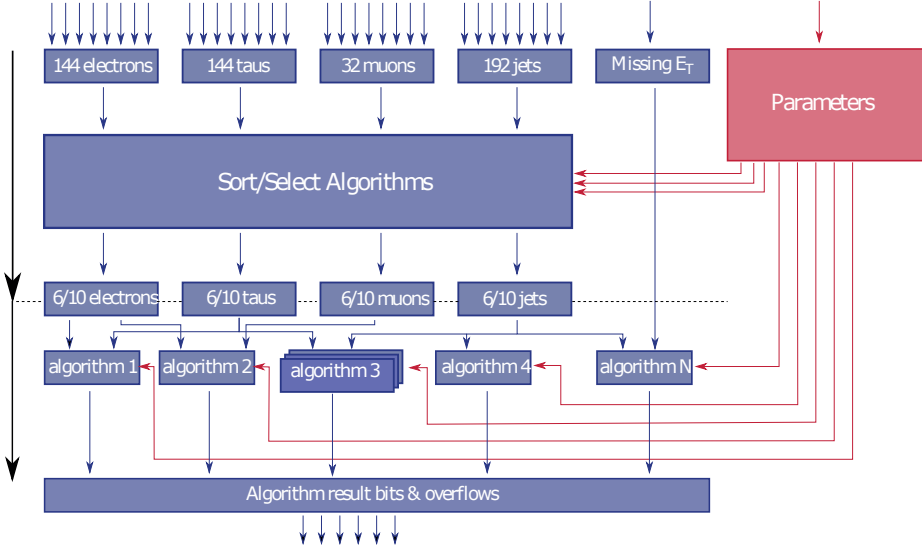
## Implementation

The KFMET algorithm is implemented on the L1Topo FPGA to be used in the Run-3 L1  $E_T^{\text{miss}}$  trigger menu. The L1Topo FPGA module, shown in Fig. 7.4, receives TOBs from the L1Calo and L1Muon modules as inputs and can combine these in more sophisticated trigger algorithms. The upgraded Phase-I L1Topo receives a  $p_T$ -sorted list containing 192 small-R jets, built out of jTowers and reconstructed at L1 with a  $0.9 \times 0.9$  ( $\eta \times \phi$ ) sliding window algorithm with a noisecut of 2 GeV [164]

The TOBs contain the  $E_T$ ,  $E_x$ , and  $E_y$  values, all with an energy resolution of 200 MeV [165]. Several hardware versus simulation tests have been conducted and show the simulation agrees with the digitized data decision made by the trigger system. These studies showed the added performance of the KFMET algorithm in the L1  $E_T^{\text{miss}}$  trigger menu and contributes towards maintaining a stable trigger rate with high efficiencies in Run-3 and beyond.

### 7.1.2 Jets without jets optimization

In order to increase the  $E_T^{\text{miss}}$  resolution on the gFEX module and reduce its dependence on  $E_T^{\text{miss}}$  originating from pile-up jets, an algorithm is designed that separates  $E_T^{\text{miss}}$  in its hard and soft components and recombines them in a linear combination. This algorithm is called the *jets-without-jets* algorithm and is optimized for identifying  $E_T^{\text{miss}}$  using L1 large-R jets built out of build out of gTowers with a  $0.8 \times 0.8$  ( $\eta \times \phi$ ) sliding window algorithm



**Figure 7.4:** Overview of the L1Topo algorithm structure. The TOB inputs are fed into Sort/Select/NoSort algorithms, which output short lists of TOBs. These lists are used and combined as inputs into various decision algorithms and the decisions are sent to the CTP [166].

[164]. The  $E_T^{\text{miss}}$  components are calculated by summing the transverse momentum in gTowers which are comprised out of supercells. The gTowers have a granularity of  $0.2 \times 0.2$  ( $\eta \times \phi$ ) in the central  $|\eta| < 2.9$  region and are combined in  $0.6 \times 0.6$  blocks in the forward region.

The hard contribution to the total  $E_T^{\text{miss}}$ ,  $H_{x,y}^{\text{gTower}}$ , is composed of the gTowers with  $p_T > 25$  GeV. The soft contribution,  $E_{x,y}^{\text{gTower}}$ , is built of all gTowers that do not pass the block cut. The total  $E_T^{\text{miss}}$  is the sum of the hard and soft contributions

$$E_{x,y,\text{jwoj}}^{\text{miss}} = a \cdot H_{x,y}^{\text{gTower}} + b \cdot E_{x,y}^{\text{gTower}},$$

where the coefficients  $a$  and  $b$  are optimized for the resolution using the Kalman Filter described in § 7.1.1 and found to be 0.98 and 0.4, respectively. This parameter combination weighs the hard contribution to the  $E_T^{\text{miss}}$  higher than the soft contribution, resulting in a decreased dependence on potential fake  $E_T^{\text{miss}}$  caused by low- $p_T$  effects. The parameters are found using the minimization target  $Q^2$ , defined as

$$Q^2 = \sum_{i=1}^N \left[ \frac{E_{x,\text{miss}}^{\text{offline}}(i) - [a \cdot H_x^{\text{gTower}}(p_T^{\text{cut}}) + b \cdot E_x^{\text{gTower}}(p_T^{\text{cut}})]}{\sigma_i} \right]^2 + \sum_{i=1}^N \left[ \frac{E_{y,\text{miss}}^{\text{offline}}(i) - [a \cdot H_y^{\text{gTower}}(p_T^{\text{cut}}) + b \cdot E_y^{\text{gTower}}(p_T^{\text{cut}})]}{\sigma_i} \right]^2,$$

which corresponds to the resolution of the algorithm using  $\sigma_i = 1$  and the same MC

sample as for the KFMET algorithm. Changing the gTower cells used in the forward region is also investigated, but does not yield any additional signal efficiency. After several iterations of algorithm parameter combinations, the setup presented in this thesis is found to be the best performing.

## 7.2 Data and Monte Carlo samples used for trigger rate estimation

At the time of the algorithm optimizations, no Monte Carlo or data-based rate estimation was available for the Run-3 L1  $E_T^{\text{miss}}$  trigger algorithms. This section briefly explains the methods and samples used to calculate a preliminary estimate of the L1  $E_T^{\text{miss}}$  trigger rates.

The studies are intended to assess the performance of  $E_T^{\text{miss}}$  algorithms, while maintaining a stable trigger rate of about 4 kHz. The data used in this study originates from an enhanced bias run [167] during the 2015-2018 data-taking period, with a  $\sqrt{s} = 13$  TeV and consists of  $57.10 \text{ pb}^{-1}$  with  $46 \leq \mu \leq 51$ . An enhanced bias dataset consists of weighted high- $p_T$  events which are likely to be selected by the trigger. It is collected using a variety of L1 triggers of all signature types and has the statistical power to assess the trigger rate of algorithmic selections on the data.

The estimated event rate is extrapolated towards Run-3 conditions by assuming a linear dependency on instantaneous luminosity [168]. To describe the background of the  $E_T^{\text{miss}}$  trigger, it is assumed to be dominated by multijet final states originating from QCD processes, particularly inelastic  $pp$ -collisions. PYTHIA 8 [110] and the ATLAS A14 [113] parton shower tune are used to model the multijet events, incorporating matrix elements calculated at LO.

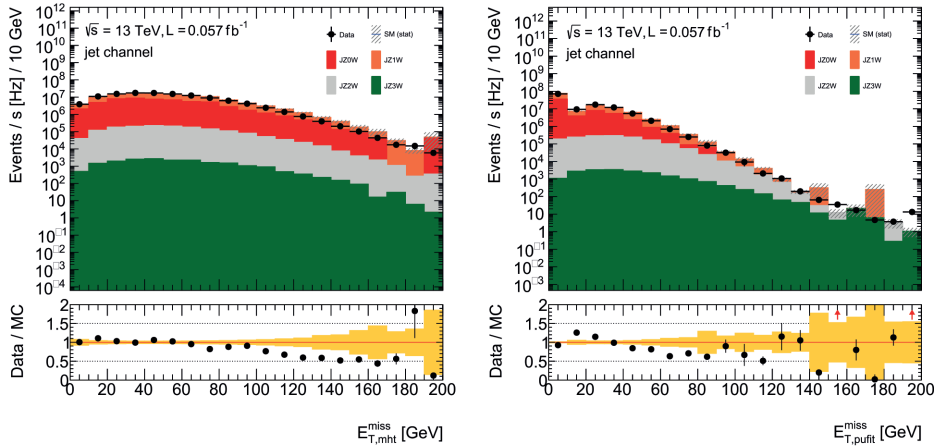
To populate the high  $p_T$  part of the inclusive jet spectrum, the sample employs a biased phase-space sampling, which is compensated by a decreasing weight for the event. On top of this, several different ranges in the  $p_T$  of the leading jet are generated separately and later combined. Since it is assumed that the background is fully dominated by inelastic di-jet scattering with low-to-medium leading jet  $p_T$ , the first four slices of jet  $p_T$  are used and shown along with their  $p_T$  range and cross-section in Tab. 7.1.

**Table 7.1:** Di-jet Monte Carlo slices with their corresponding leading jet  $p_T$  range and cross-section.

| Di-jet slice | Leading jet $p_T$ [GeV] | Cross-section [pb]    |
|--------------|-------------------------|-----------------------|
| JZ0          | 0–20                    | $7.81 \times 10^{10}$ |
| JZ1          | 20–80                   | $7.84 \times 10^{10}$ |
| JZ2          | 80–200                  | $2.43 \times 10^9$    |
| JZ3          | 200–500                 | $2.65 \times 10^7$    |

### 7.3 The trigger performance and rate analysis

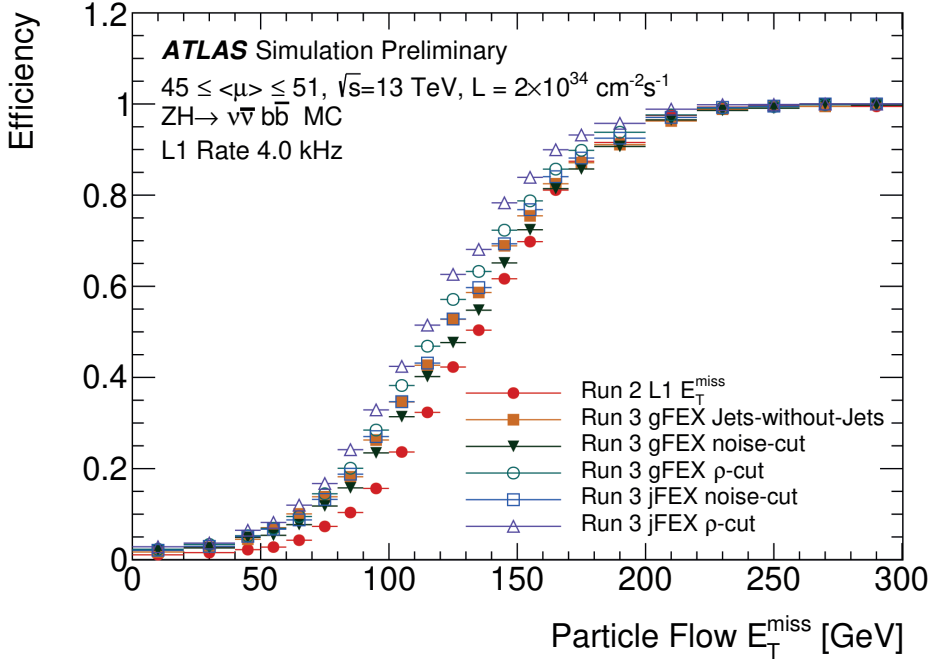
In this section, the di-jet simulation is compared to the emulated Run-3 enhanced bias dataset. At least one L1-jFEX jet with a radius parameter  $R = 0.4$  and  $p_T$  of 5 GeV is required. The events are rescaled to the total number of events per second (Hz) in each bin by using an event-by-event weight. The resulting rate distributions of the HLT  $E_T^{\text{miss}}$  algorithms are depicted in Fig. 7.5 and show the unbiased pass rate of two HLT  $E_T^{\text{miss}}$  algorithms using events with  $46 \leq \mu \leq 51$ .



**Figure 7.5:** The HLT  $E_T^{\text{miss}}$  mht rate (left) and the HLT pufit  $E_T^{\text{miss}}$  rate (right). The JZ0W, JZ1W, JZ2W and JZ3W processes correspond to the di-jet slices in Tab. 7.1.

Although the di-jet Monte Carlo seems to over-predict the rate at  $E_T^{\text{miss}} \sim 50\text{--}100$  GeV, this agreement was deemed sufficient for use in the prediction of the integrated rate for  $E_T^{\text{miss}}$  trigger algorithms.

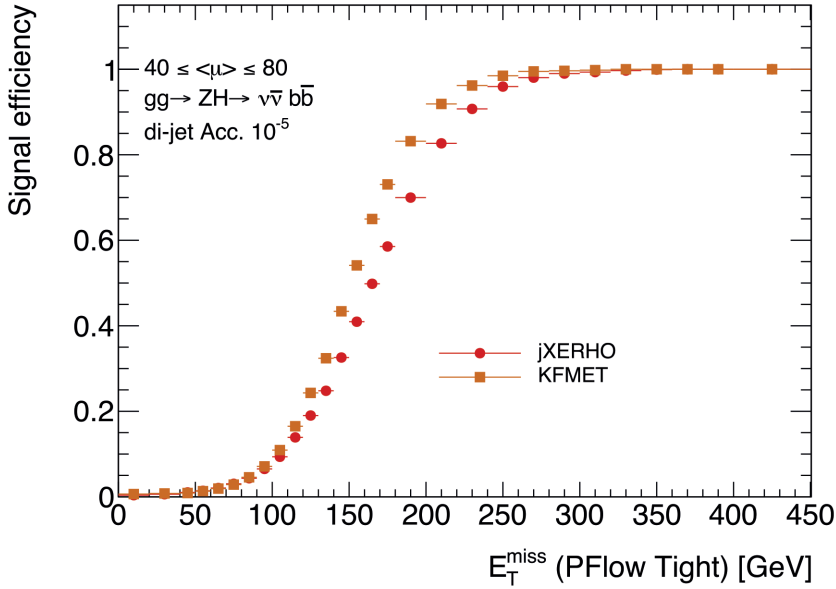
In Fig. 7.6, the rate prediction is integrated across the predicted  $E_T^{\text{miss}}$  and used to predict and compare the efficiencies of the proposed gFEX and jFEX  $E_T^{\text{miss}}$  algorithms while maintaining equal trigger rate, determined by data, on a  $ZH$  signal topology. By using the equal rate efficiency comparison, it is shown that the best performing Run-3 L1Calo algorithm is outperforming the legacy Run-2  $E_T^{\text{miss}}$  trigger by  $\mathcal{O}(15\%)$  at an offline  $E_T^{\text{miss}}$  of 150 GeV. This allows for operation at equal efficiency while greatly decreasing the total trigger rate; or operate at higher efficiency while maintaining Run-2 trigger rates.



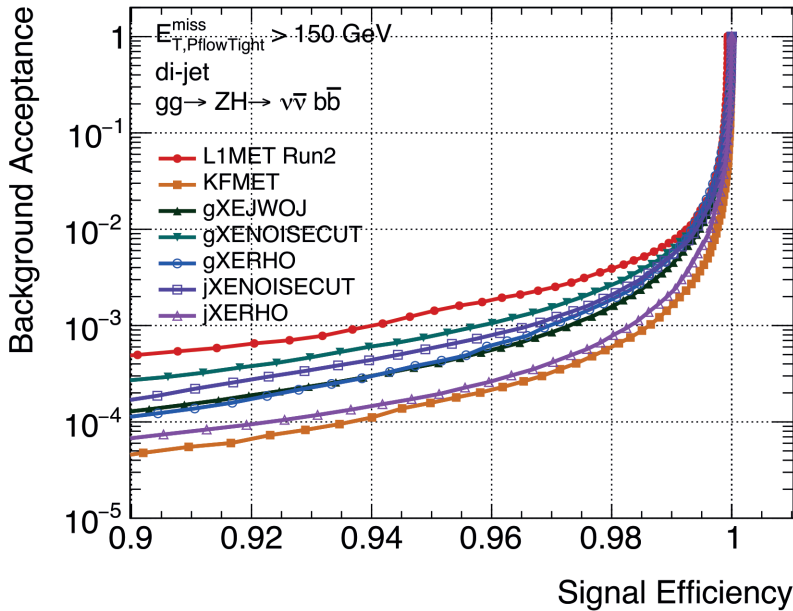
**Figure 7.6:** Equal rate efficiency of the new missing transverse momentum algorithms proposed for the Run-3 jFEX and gFEX, compared to the Run-2  $E_T^{\text{miss}}$  trigger. The efficiency is computed from  $ZH \rightarrow \nu\bar{\nu}b\bar{b}$  Monte Carlo simulation with respect to the offline  $E_T^{\text{miss}}$  using the *Tight* working point. All thresholds are tuned to give approximately the same rate as the Run-2 trigger. The *noise-cut* algorithm computes  $E_T^{\text{miss}}$  from the vector sum of all towers with  $E_T$  above an  $\eta$ -dependent threshold. The  $\rho$ -cut algorithm computes  $E_T^{\text{miss}}$  from the vector sum of all towers with  $E_T$  above a threshold depending on  $\eta$  and the local per-event pileup density ( $\rho$ ). The jets without jets algorithm computes  $E_T^{\text{miss}}$  based on a linear combination of the soft and hard contributions to the  $E_T$  of all towers [2].

In Fig. 7.7, the efficiency of the KFMET algorithm is shown for a  $ZH$  sample as compared to the best performing algorithm at L1,  $\rho$ -cut (jXERHO) at equal rate, estimated via di-jet Monte Carlo. The KFMET reaches an  $\mathcal{O}(10\%)$  efficiency improvement at an offline  $E_T^{\text{miss}}$  of 200 GeV.

Additionally, in Fig. 7.8, the KFMET ROC curve is compared against the other proposed algorithms on the gFEX and jFEX modules. The KFMET algorithm is found to be outperforming the best jXERHO algorithm at L1 and showing the highest efficiency at equal background acceptance, defined as the ratio of accepted and total background events. The ROC curve shows that the KFMET algorithm also has the highest efficiency at the lowest rate, making it the best performing algorithm for the Run-3 L1  $E_T^{\text{miss}}$  trigger menu at any desired trigger efficiency or trigger rate.



**Figure 7.7:** Efficiency of the KFMET missing transverse momentum ( $E_T^{\text{miss}}$ ) algorithm for Run-3 L1Topo compared to its input base  $E_T^{\text{miss}}$   $\rho$ -cut algorithm (jXERHO) displayed for a  $ZH$  sample, evaluated at equal rate.



**Figure 7.8:** The ROC curve, derived from the  $ZH$  sample, of the KFMET algorithm is shown and compared against the other proposed algorithms on the gFEX and jFEX modules. All events are subject to an offline  $E_T^{\text{miss}} > 150$  GeV selection using the *Tight* working point. The background acceptance is defined as the ratio of accepted and total background (di-jet) events.



## THE RECONSTRUCTION OF PROTON-PROTON COLLISIONS

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When measurements are performed, the data accumulated by the ATLAS experiment and the digitized simulation go through identical reconstruction algorithms in order to be able to compare them. In this chapter, we discuss how the reconstruction of  $pp$ -collisions in the ATLAS detector is performed and describe the various reconstruction algorithms used to convert the measured signal into physical objects. In order to reconstruct the signature of interest in this thesis, the reconstruction of the following objects is essential: tracks, vertices, electrons, muons, jets and missing energy. The inclusion of taus in the signal's signature is also briefly touched upon. The reconstruction of these objects is described in the following sections.

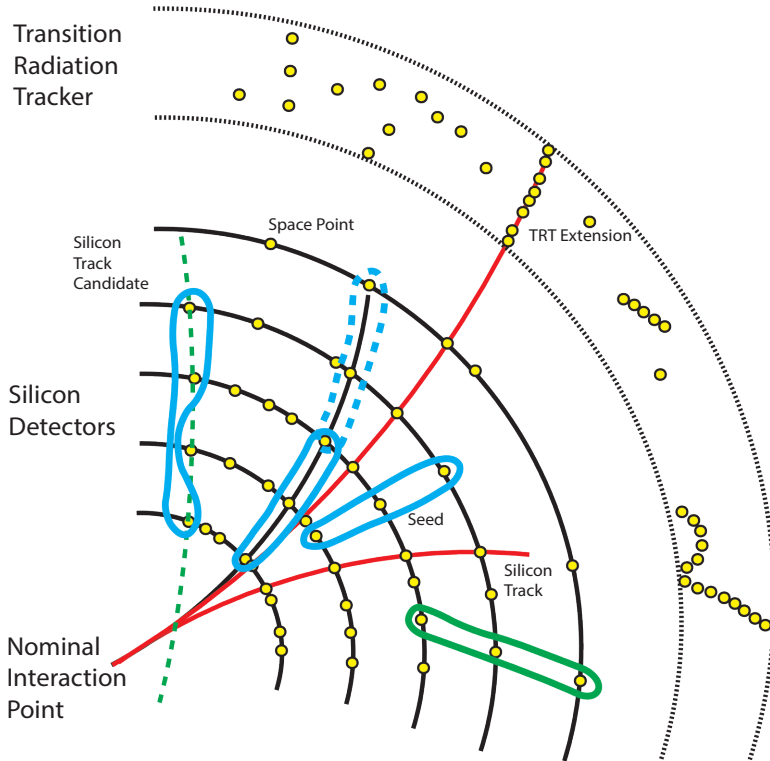
### 8.1 Tracks and vertices

Because the charged particles pass through the various layers of the ID, they leave signals in the detector, which are referred to as hits. A charged particle usually leaves multiple hits inside the detector from which a track is constructed. The hits are supplied to the track reconstruction algorithm which finds the best-fit for the trajectory of the particle. These tracks are in turn used to determine the physical trajectory the particle takes through the detector. The tracks are then used to determine the vertex of the collision, which is the point in space where the  $pp$ -collision occurred or the particle is produced.

The ATLAS track reconstruction algorithm [169, 170] receives information from the ID in the form of hits, it begins by building clusters of hits. The hits originate from the two silicon sub detectors in the ID: the Pixel Detector and the Semi-Conductor Tracker (SCT). As the hits are initiated by exceeding the energy deposition threshold of the silicon strips or sensors, they form the signal of a passing charged particle. Clusters are combined into raw neighbouring hit groups and can be interpreted as originating from the same candidate particle. Finally, the space points, which are formed using the clusters, are the 3D points in space where the hits are located; they represent the position where the charged particle traverses the active material of the ID. In the pixel detector, each cluster forms one space point, whereas in the SCT, clusters from both sides of a strip layer are combined to form a space point. This decreases the chance of misidentifying a cluster of hits as a space point.



After identification of the space points by the track reconstruction algorithm, the algorithm then finds the best-fit for the trajectory of the particle. The inputs for this step are so-called track seeds, which are formed by sets of three space points that are close to each other. A combinatorial search is then performed by a Kalman filter [171] to build track candidates from the track seeds. This process, illustrated in Fig. 8.1, is able to create multiple track candidates from the same track seeds. As the quality of space points depends heavily on the used subdetectors and the amount of material the particle traversed, different seed types are considered.



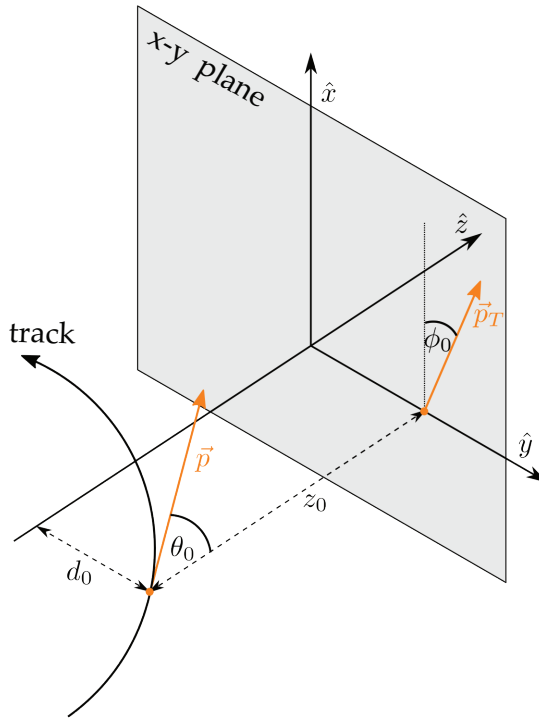
**Figure 8.1:** Illustration of the stages of the track pattern recognition in a simplified model. The space points are shown in yellow. The track seeds reconstructed from combinations of space points are outlined in blue. The dashed blue seed illustrates a case where two seeds correspond to the trajectory of the same charged particle. The green seed is rejected by the requirement that the resulting track should be consistent with the nominal interaction point. The green dashed line is a track candidate, rejected because it is inconsistent with the nominal interaction point. The black track candidate is a fully reconstructed silicon track with no TRT extension. A track reconstructed using information from all three subdetectors is shown in red [172].

First, SCT-only seeds are used, followed by Pixel-only seeds and finally SCT-Pixel seeds. This method results in a set of track candidates that are then filtered by a track fitter which selects the most likely track candidate; this is referred to as the ambiguity-solving stage. The track fitter uses a  $\chi^2$  minimization technique and various track scores to determine the best-fit for the track; criteria like the number of hits, the  $\chi^2$  of the fit, the

number of holes in the track, the number of shared hits and the calculated momentum are used to determine the best track candidate.

The track fitter limits the number of track candidates that share clusters, gives priority to tracks established in the Pixel Detector and retains efficiency on the track building using merged clusters and a neural network for pixel clustering. At last, the surviving track candidates are extended into the TRT, which has a worse resolution compared to the SCT and Pixel Detector. The TRT is used to extrapolate the track further into the detector and improve the momentum resolution.

It is assumed the resulting track can be described with a perfect helical trajectory in a uniform magnetic field, depicted in Fig. 8.2. The tracks are fitted using the information from all three layers of the ID and are described by five parameters, being respectively: the charge over momentum,  $q/p$ , the polar angle,  $\theta$ , the azimuthal angle,  $\phi$ , the impact parameter,  $d_0$ , and the  $z$ -coordinate of the longitudinal difference along the beam line between the point where the impact parameter is measured and the primary vertex,  $z_0$ . The track reconstruction algorithm is able to reconstruct tracks with a transverse momentum  $p_T$  of at least 400 MeV and a pseudorapidity of  $|\eta| < 2.5$ . The impact parameter,  $d_0$ , is required to be less than 2 mm and the longitudinal impact parameter,  $z_0 \times \sin\theta$ , is required to be less than 3 mm.



**Figure 8.2:** A track candidate with a helical shape described by the transverse and longitudinal track parameters  $d_0$  and  $z_0$  respectively. Along with the polar angle  $\theta_0$ , the azimuthal angle  $\phi_0$ , momentum  $\vec{p}$  and transverse momentum  $\vec{p}_T$ . Taken from Ref. [173].

The vertex associated to a subset of charged particle tracks is found by a dedicated vertex finding reconstruction algorithm [174, 175]. An accurate reconstruction of the number and position of interaction vertices is needed to separate the pile-up interactions from the hard-scattering vertex. Additionally, b/c-hadron decays are displaced and produce a secondary vertex which is used in the determination of the flavour of the quark that initiated the jet. Therefore, the vertex finding is essential for the physics analysis being performed using the ATLAS detector.

The primary vertex reconstruction happens in two stages: the first stage is the vertex finding, and the second stage is vertex fitting. The first stage is the pattern recognition process, in which reconstructed tracks are associated with vertex candidates. The vertex candidates are formed by selecting tracks that are consistent with originating from the same point in space and they are subject to quality criteria on the impact parameters and tracks must have  $p_T > 500$  MeV. After the candidates are formed, the iterative vertex fitting procedure is started, which is a  $\chi^2$  minimization technique used to determine the best-fit for the vertex. The seed position of the vertex fit is based on the collision point in the transverse plane, where the  $xy$ -coordinates of the starting point are taken from the centre of the collision point, which is reconstructed using a dedicated beam spot reconstruction likelihood fit. The  $z$ -coordinate of the starting point is calculated as the mean  $z$ -coordinate of the tracks at their respective points of closest approach to the reconstructed centre of the collision point.

After candidates are established, the vertex fitting procedure is started, in which the best position of the nominal vertex is found. This is achieved by assigning each input track a weight reflecting the compatibility of the track with the vertex, then the vertex position is recalculated using these reweighted tracks. In this calculation, tracks with a relatively large weight have a larger influence on the vertex position. After the new reweighted vertex position is found, the track weights are recalculated using the new vertex position. This process is repeated until the vertex position converges to a stable position. After this procedure, tracks with a sufficiently large weight remain assigned to the vertex and tracks with a low weight are removed from the vertex, which are later used for the determination of the next vertex in the ID, until no new vertex can be found. The vertex whose tracks has the highest  $\sum p_T^2$  is assumed to be the hard-scatter vertex and the other vertices are considered pile-up vertices.

## 8.2 Calorimeter measurements

Now tracks and vertices are identified using the ID, the information is combined with the energy measurements of the EM and hadronic calorimeters within the ATLAS detector. By combining the tracking and energy measurement elements, all the main physical objects are reconstructed.

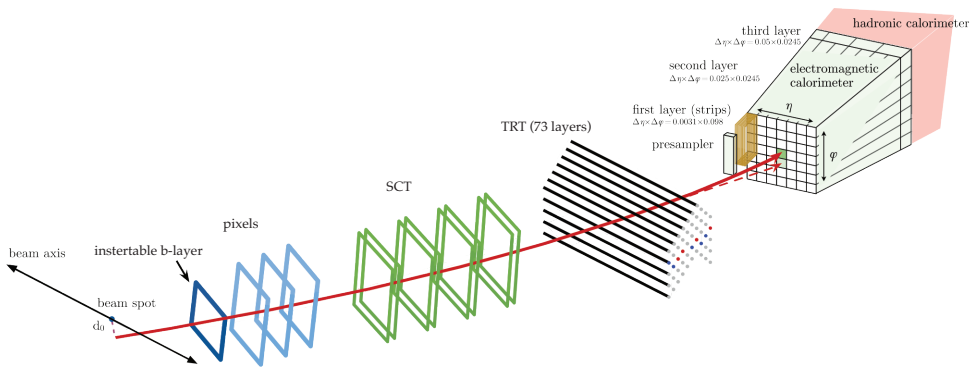
The energy deposited by a particle shower is captured by the hadronic and electromagnetic calorimeters and is first described in the form of topological clusters [176, 177]. The

algorithm begins by forming clusters of calorimeter cells in both hadronic and electromagnetic calorimeters that have readouts above a certain energy threshold. The clusters can be viewed as three-dimensional energy deposits of particle showers inside the active material of the calorimeters. The calorimeter cell initiating the cluster is required to have an energy  $|E_{\text{cell}}| > 4\sigma_{\text{noise}}$ , where  $\sigma_{\text{noise}}$  is the average energy noise level of a cell, in order to distinguish from the background noise inside the calorimeters. Additional neighbouring cells are added to the cluster if they have an energy  $|E_{\text{neighbouring cell}}| > 2\sigma_{\text{noise}}$ .

This process is repeated until no more cells can be added to the cluster, after which a final round of cells surrounding the cluster, satisfying  $|E_{\text{neighbouring cell}}| > 0$ , is added to the cluster to account for the energy leakage from the shower. At last, the resulting cluster is checked for maxima exceeding 500 MeV. If more than one is found, the cluster is split such that only one maximum is present in each cluster which represents a particle per cluster.

### 8.2.1 Electrons

By combining calorimeter energy cluster measurements in the ECal and tracks in the ID, electrons are reconstructed using a dedicated reconstruction algorithm [177, 178]. Fig. 8.3 depicts an electron and a radiated photon travelling from the beam spot towards the calorimeters, illustrating the various subdetectors which deliver information to the reconstruction of electrons.



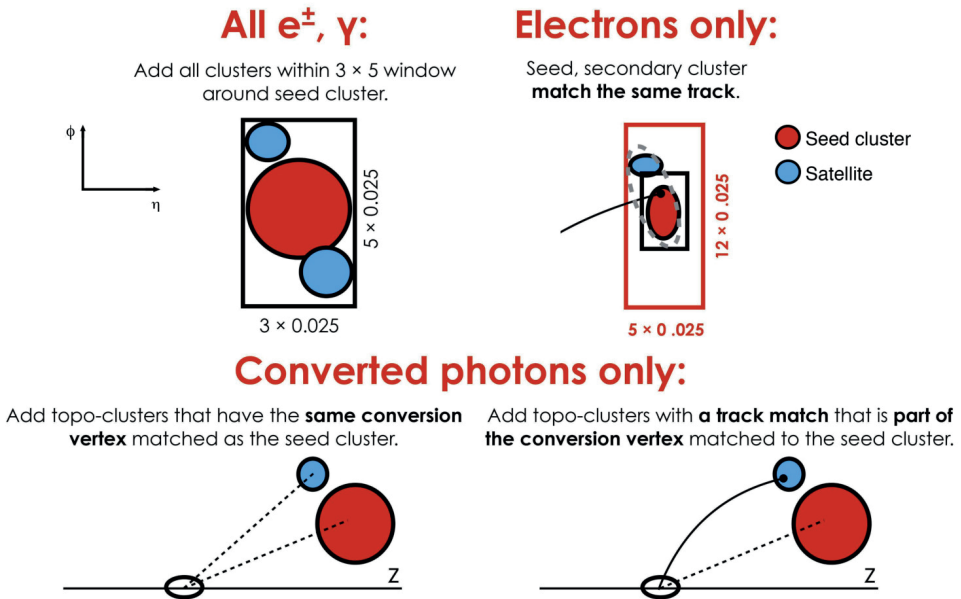
**Figure 8.3:** A schematic illustration of the path of an electron through the detector. The red trajectory shows the hypothetical path of an electron, which first traverses the tracking system (pixel detectors, then silicon-strip detectors and lastly the TRT) and then enters the electromagnetic calorimeter. The dashed red trajectory indicates the path of a photon produced by the interaction of the electron with the material in the tracking system. Adapted from [177].

When electrons interact with the material they traverse, they lose a significant amount of their energy due to bremsstrahlung, radiating photons which can further convert into electron-positron pairs. This results in a shower of particles which is detected by the ECal. Despite the fact photons leave no track in the ID, they are similarly detected and reconstructed, but as electromagnetically neutral particles using only the energy clusters

in the ECal.

The electron and photon reconstruction uses a dynamic variable-size clustering to improve the energy resolution, which is deteriorated by photons radiating off electrons and charged particles from pile-up. The clustering into superclusters starts with an incoming list of EM topological clusters sorted according to descending  $E_T$ . For a cluster to become a seed cluster, it is required to have at least an  $E_T$  of 1 GeV and to be matched to a track with at least four hits in the silicon track detectors. For photons, the  $E_T$  is required to be at least 1.5 GeV, without a matched track.

When a seed cluster is identified, the supercluster algorithm attempts to find satellite clusters close to the seed cluster, using the process illustrated in Fig. 8.4. For  $e^\pm$  and  $\gamma$ , clusters are labelled a satellite cluster when it is found within a window of  $\Delta\eta \times \Delta\phi = 0.075 \times 0.125$  around the seed cluster barycentre. The barycentre is the weighted average of the supercluster centre in the  $\eta - \phi$  plane. Additionally for  $e^\pm$ , clusters are also considered within a window of  $\Delta\eta \times \Delta\phi = 0.125 \times 0.300$  around the seed cluster barycentre that share their best matched track with the seed cluster. For  $\gamma$  with conversions vertices, the satellite best-matched tracks are required to belong to the same conversion vertex as the seed cluster.

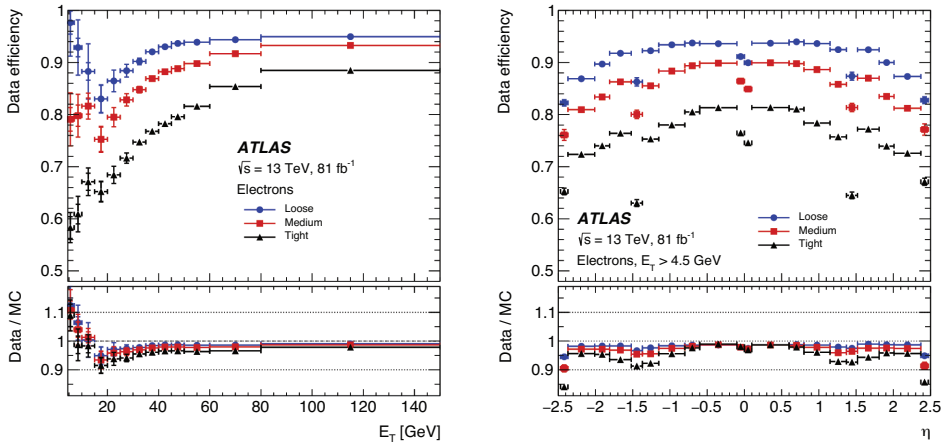


**Figure 8.4:** Visualization of the superclustering algorithm for electrons and photons. Seed clusters are shown in red, satellite clusters in blue [178].

To improve the reconstructed momentum and energy, the initial energy calibration and position corrections are estimated for the superclusters by comparing simulated  $Z \rightarrow ee$  and  $J/\psi \rightarrow ee$  events with data. During the calibration, the tracks are matched to the electron superclusters and converted vertices to the photon superclusters. From

these calibrated superclusters, the final electron and photon candidates are reconstructed. The final calibration is performed using boosted decision trees, using multiple variables representing the energy fractions in different layers of the calorimeter, the shower shape, the track-cluster matching and the track quality [179].

At last, the resulting electrons can be classified on the basis of their reconstruction quality into an identification classification. In order to distinguish their signature from photons and jets, the identification and isolation criteria make use of the activity around the object to suppress potential backgrounds. These are applied in the form of available working points called *VeryLoose*, *Loose*, *Medium* and *Tight* and are shown in Fig. 8.5.



**Figure 8.5:** The electron identification efficiency in  $Z \rightarrow ee$  events in data as a function of  $E_T$  (left) and as a function of  $\eta$  (right) for the *Loose*, *Medium* and *Tight* operating points. The efficiencies are obtained by applying data-to-simulation efficiency ratios measured in  $J/\psi \rightarrow ee$  and  $Z \rightarrow ee$  events to  $Z \rightarrow ee$  simulation. The inner uncertainties are statistical and the total uncertainties are the statistical and systematic uncertainties in the data-to-simulation efficiency ratio added in quadrature. For both plots, the bottom panel shows the data-to-simulation ratios [179].

The looser selection benefits from high efficiency but lower purity, while the tighter selection has a higher purity but lower efficiency. The same approach holds for photons.

In order to reduce the number of mis-identified leptons and acceptance of non-prompt leptons, which are one of the major backgrounds in  $H \rightarrow WW^*$  analyses and mostly produced through the production of a  $W$  boson in association with a jet, additional criteria on the isolation of electrons are required. The newly introduced *PLIVImprovedTight* [180], the loosest of the two improved prompt lepton-veto working points, reduces the number of mis-identified leptons by a factor of 2 compared to the *Tight* isolation working points. Therefore, the new working points are used in the measurements in this thesis which analyze the full Run-2 dataset of  $140 \text{ fb}^{-1}$ .

The PLIV working point introduces many new improvements compared to their older isolation criteria, one of the major is the usage of a recurrent neural network trained in

Keras [181] to distinguish prompt leptons from non-prompt leptons. The neural net makes use of additional isolation variables, such as the collection of hadron energy contributions around the calorimeter clusters. The working point uses a dedicated secondary vertex reconstruction algorithm to identify the non-prompt leptons produced in the b/c-hadron decays. In this algorithm, the reconstructed secondary vertices are used to estimate the flight path of a hadron before it decays to a non-prompt lepton.

Although the signal efficiency is reduced by up to 10%, it is beneficial for the analyses to use *PLIVImprovedTight* as the signal significance increases; it yields a 2% increase in the Poissonian counting significance compared to the stricter *PLIVImprovedVeryTight* working point. Utilising the *PLIVImprovedTight* working point additionally benefits the quality of the modelling in validation regions enhanced with mis-identified leptons, such as events with two leptons with the same reconstructed electric charge.

### 8.2.2 Muons

The muons produced in the  $H^* \rightarrow WW$  decay typically make their way through the ID and arrive at the muon spectrometer. To reconstruct the muons, the algorithm [182, 183] uses the tracks in the ID and the hits in the muon spectrometer.

First, the MS track reconstruction searches for hits in the muon chambers to reconstruct the muon track. Hits that follow a bending trajectory expected for a charged particle in a magnetic field are combined into segments. Compatible segments in each of the muon chambers in the MS are combined and fitted into a stand-alone muon track. The building of a muon track requires at least two segments to be matched to each other, based on their relative position and angles, except for the MS transition region where only one is needed. After an overlap removal algorithm addresses ambiguity in the muon reconstruction by removing muon tracks candidates that have a shared segment, a  $\chi^2$  fit is used to select the final muon track candidate.

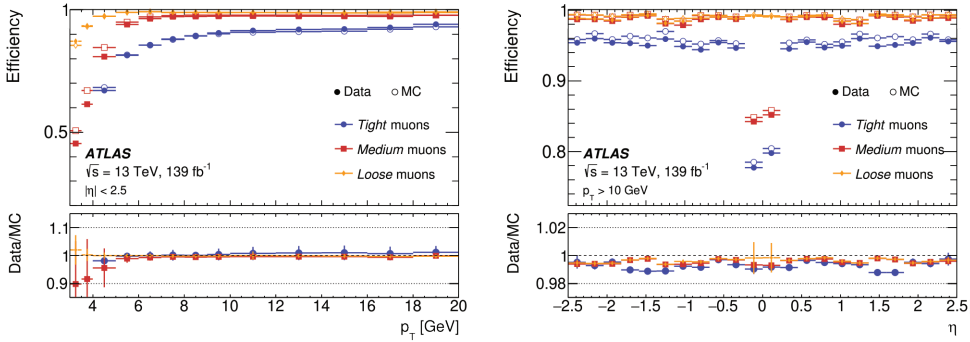
After assembly, the ID and MS track candidate calorimeter information is added and the combined muon reconstruction is performed. Various algorithms are used in this reconstruction that produce four muon types, based on the which subdetectors used in the reconstruction.

- **Combined muons** are reconstructed using both the ID and MS tracks and are fitted to make a combined track which results in reconstructed muons with the least amount of systematic uncertainty. During this fit, MS hits may be added or removed to improve the fit quality. In this algorithm, the outside-in method is used as default, where MS tracks are extrapolated inwards. If the reconstruction fails, an inside-out method can also be used which extrapolates ID tracks towards the MS.
- **Segment-tagged muons** are reconstructed using a track in the ID which is matched to at least one track segment in the MDT or CSC chambers. This is either because they have low- $p_T$  or fall in regions with reduced acceptance.



- **Calorimeter-tagged muons** are constructed from a track in the ID which is matched to a calorimeter energy cluster compatible with a minimum-ionising particle. These muons are used to recover acceptance in the regions where the ATLAS MS has decreased acceptance which is at  $|\eta| < 0.1$ .
- **Extrapolated muons** are reconstructed using only MS tracks and a loose compatibility requirement to originate from the interaction point. These muons are used to recover acceptance in the regions where the ATLAS ID is not present, which is at  $2.5 < |\eta| < 2.7$ .

After the muon candidates have been established, they are calibrated in order to improve their energy and momentum reconstruction. This is achieved by comparing  $Z \rightarrow \mu\mu$  and  $J/\psi \rightarrow \mu\mu$  data against simulation, as shown in Fig. 8.6.



**Figure 8.6:** Muon reconstruction and identification efficiencies for the *Loose*, *Medium*, and *Tight* criteria. The left plot shows the efficiencies measured in  $J/\psi \rightarrow \mu\mu$  events as function of  $p_T$  with  $|\eta| < 2.5$ . The right plot displays the efficiencies measured in  $Z \rightarrow \mu\mu$  events as a function of  $\eta$ , for muons with  $p_T > 10$  GeV. The predicted efficiencies are depicted as open markers, while filled markers illustrate the result of the measurement. The statistical uncertainty in the efficiency measurement is smaller than the size of the markers, and thus not displayed. The panel at the bottom shows the ratio of the measured to predicted efficiencies, with statistical and systematic uncertainties [183].

Similar to the electron and photon reconstruction, muons are also separated into three main identification categories: these categories are *Loose*, *Medium* and *Tight*. Also similar to electrons, these working points balance signal efficiency against purity, allowing the analyser to choose their preferred working point. In addition to the identification criteria, the muons are also required to pass isolation criteria. The isolation criteria are used to suppress the acceptance of non-prompt muons or the misidentification of jets as muons. In this thesis, the *FCTight* isolation working point is used along with the new *PLIVImprovedTight* [180] identification criteria.



### 8.2.3 $\tau$ leptons

The  $\tau$  lepton is the heaviest lepton, with a mass of 1.777 GeV [184]. It decays either leptonically, in a one-prong decay, into a less heavy electron or muon, along with two neutrinos, via the weak charged boson decay  $\tau \rightarrow W\nu_\tau \rightarrow \ell\nu_\ell\nu_\tau$ ,  $\ell = e, \mu$ ; it can also decay hadronically, using a one-prong decay into a  $\pi^\pm$  and a  $\nu_\tau$ , leaving one track, or a three-prong decay into  $\pi^\pm\pi^\pm\pi^\mp$  and a  $\nu_\tau$ , leaving three tracks. Both decays typically take place before any active parts of the ATLAS detector are reached. In this thesis, the reconstructed  $WW \rightarrow \ell\nu\ell\nu$  signatures includes potential leptonic  $\tau$  decays which enter the analysis regions via their decay products as electrons or muons. The ATLAS reconstruction also features dedicated reconstruction of hadronic  $\tau$  decays [185], but these are not targeted by measurements presented in this thesis; the hadronic  $\tau$  decays are accounted for when performing overlap removal for the reconstruction of different objects.

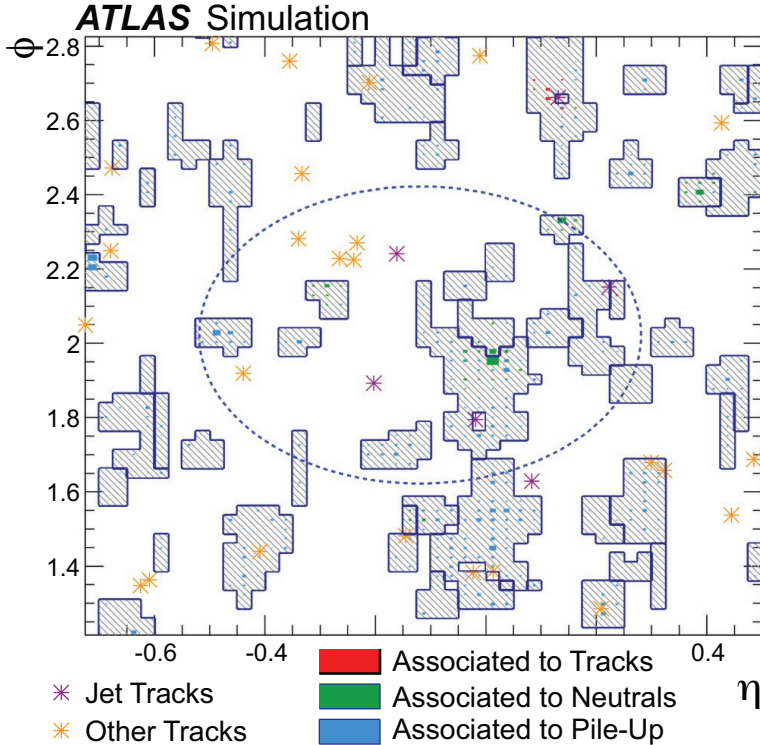
## 8.3 Jets

Jets are the most common final-state objects in  $pp$ -collisions and are often produced in association with the Higgs signal of interest in this thesis. Both the ggF and VBF Higgs boson production modes are characterized by the potential presence of jets in the final state. While for VBF the presence of two forward jets is a characteristic signature, for ggF the presence of jets is due to the ISR originating from the gluon-initiated top loop. The jet characteristics are also used to distinguish the signal from the background, to normalize the background in various jet multiplicity regions and it plays a large role in higher-order QCD corrections to the gluon-initiated Higgs production.

The nature of QCD, and its colour-charge, causes partons produced in the collision and those radiated by the initial state to either form additional partons or to hadronize into a collimated shower of particles. The showers of particles are detected from the tracks and energy deposits and are then reconstructed as jets. The jets are reconstructed according to the Particle Flow (PFlow) method [186]: this uses tracks in the jet reconstruction to improve the resolution of the jets, it allows for better pile-up rejection and improves the angular resolution. The tracks are provided by the ID and thus only cover the central region of the detector  $|\eta| < 2.5$ .

The first step of the jet reconstruction is to remove overlaps between the reconstructed input objects. This is done by removing the energy deposits in the calorimeter that are associated with, for example, reconstructed electrons, photons or muons. By matching well-measured tracks with energy deposits in the topoclusters, additional topoclusters are added since particles can deposit energy multiple clusters. At this point, the expected energy deposited in the calorimeter is subtracted, cell by cell, from the set of matched topoclusters.

Following this, the tracking and calorimeter information is combined to form the jets. This is illustrated in Fig. 8.7, where the second layer of the EM calorimeter is shown. A 30 GeV truth jet is outlined with an ellipse of  $R = 0.4$ , the radius of the jet in the  $\eta$ - $\phi$  plane of the detector. Two types of jets are defined: hadronic jets and soft-activity jets. Jets defined as soft activity fail the  $p_T$  threshold used in the jet reconstruction and are reconstructed separately for later usage.



**Figure 8.7:** A display of  $R = 0.4$  jet reconstruction in the second layer of the EM calorimeter with a  $p_T$  of 30 GeV. Asterisks indicate the position of tracks extrapolated to the calorimeter: a purple asterisk indicates tracks to be used for particle flow jet reconstruction, while other tracks are shown in orange asterisks. The red and green boxes indicate reconstructed particles where tracks have not been used and blue boxes are associated to pile-up. The blue diagonally striped cells are clustered into topoclusters [186].

Crucial in the jet reconstruction is the usage of a jet-finding algorithm. As jets are not fundamental physical objects, but a definition of a shower of particles, the definition of a jet can differ from experiment to experiment. Additionally, the jet-finding algorithm must be consistent between real  $pp$ -collisions and those from a Monte Carlo simulator. The jet-finding algorithm used in ATLAS is the anti- $k_t$  algorithm [187]: this algorithm is a sequential recombination algorithm which defines how to group the particles into jets, be it simulated particles or reconstructed particles.

Another requirement for the jet-finding algorithm is that it should be IR and collinear safe. This means that, in the case of collinear safety, the jet algorithm and its resulting jets should not be sensitive to the splitting of one particle into two softer collinear particles carrying the same four momentum. In the case of IR-safety, the definition of the jet should be stable in the presence of low- $p_T$  radiation of partons in the event. The algorithm takes objects  $i$  and  $j$ , which can be partons or reconstructed particles, and calculates the distance between them, which is calculated from their transverse momentum  $k_t$ , the rapidity and the azimuthal angle:

$$d_{ij} = \min(k_{t,i}^{-2}, k_{t,j}^{-2}) \frac{\Delta_{ij}^2}{R^2},$$

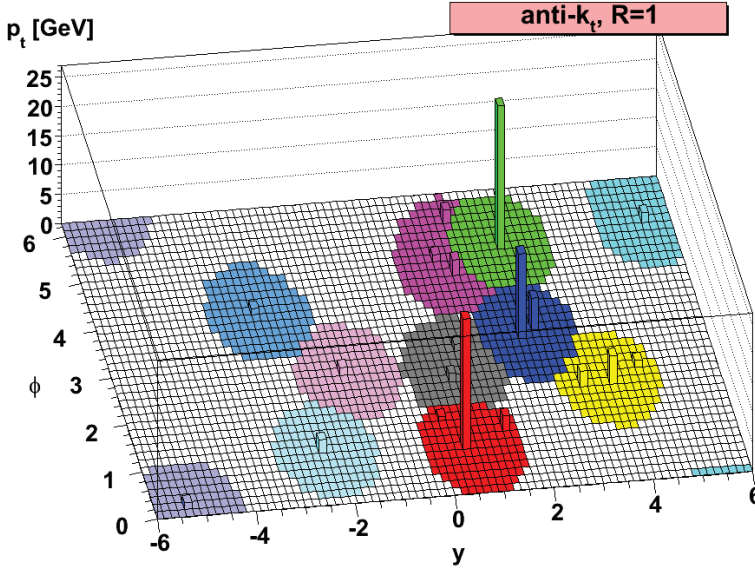
$$d_{iB} = k_{t,i}^{-2}.$$

In this equation, the distance between particles is defined by  $\Delta_{ij}^2 = (y_i - y_j)^2 + (\phi_i - \phi_j)^2$ , the radius of the jets is  $R$ . The starting point for the distances is  $d_{iB}$ , which denotes the distance between object  $i$  and the beam  $B$ . Mostly small- $R$  jets ( $R = 0.4$ ) are used in this thesis. However,  $R = 1$  jets, which are referred to as large- $R$  jets, are often used when describing objects with large  $p_T$ .

Because the clustering algorithm relies on the  $k_t$  of a particle, soft particles tend to cluster with hard particles before they cluster among themselves. Therefore, if a hard particle has no hard neighbours within a distance of  $2R$ , it accumulates all the soft particles and result in a perfect cone jet originating from the hard particle. However, in the case of two hard jets, with  $k_{t,1} \gg k_{t,2}$  and outside of each other's cone radius, the first jet is of conical shape, while the second jet misses the overlap with jet 1. For a third situation,  $k_{t,1} \approx k_{t,2}$ , both cones are clipped, with the boundary  $b$  between them defined by  $\Delta_{R1b}/k_{t,1} = \Delta_{2b}/k_{t,2}$ . For a fourth situation with,  $\Delta_{12} < R$ , the two particles are combined into a single jets and the shape becomes more irregular, based on the momentum ratio of the two hard particles.

These various situations are illustrated in Fig. 8.8, where a hard-scatter parton-level event is shown with random soft ghost particles added to the event. This illustrates the ability of the algorithm to form cone-shaped jets, irrespective of the ordering of the soft particles.

Now the jets have been clustered and reconstructed, they are calibrated. This is performed in a series of steps that correct for detector effects and alter the original  $E_T$ ,  $p_T$  and mass of the particles and the resulting jets [189]. The first step is the Origin Correction, which corrects the jet four-momentum to point to the primary vertex instead of the centre of the detector. Due to the large amount of in-time and out of time pileup during collisions, a dedicated pile-up correction method is developed. First a per-jet area pile-up subtraction is performed, which subtracts the median energy density,  $\rho = p_T^{\text{jet}}/A$ , from the jet energy. The area  $A$  in the median energy density definition is the area in which the jet deposits energy. After this, another global residual pile-up correction is performed, based on the



**Figure 8.8:** A hard-scatter parton-level event generated with Herwig and random soft ghost particles added to the event. The anti- $k_t$  algorithm is used to cluster the particles into jets and are shown as cones with a radius of  $R = 1$  [188].

ratio of the  $\rho$ -subtracted jet  $p_T$  to the uncorrected jet  $p_T$ . This is taken as a correction factor to the jet four-momentum which does not affect the jet  $\eta$  and  $\phi$  coordinates.

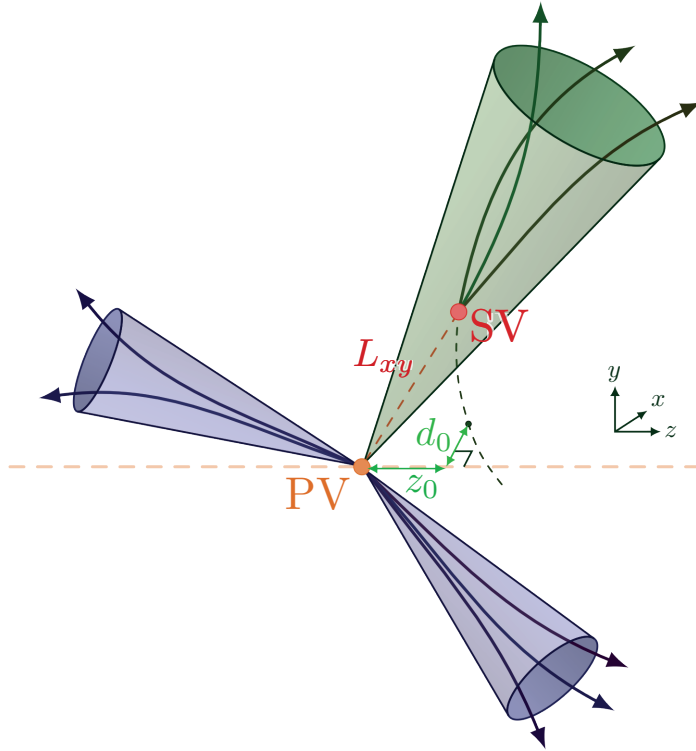
The final step is the Jet Energy Scale (JES) calibration, which corrects the jet's four momenta to the particle-level using a MC-based calibration. This is achieved by comparing the reconstructed jets to the truth jets in the simulation and correcting for the non-linear response of the calorimeters. The global sequential calibration reduces dependence on the jet being quark or gluon initiated, leading to a more uniform response for the jets. The last in-situ calibration method corrects for the difference in jet response between data and simulation that arises from the imperfect detector description of the jet response and detector material in the MC simulation.

In order to further improve the reconstruction of jets, a Jet Vertex Tagger (JVT) [190, 191] is employed. This tagger uses a multivariate combination of the variables  $R_{p_T}$  and CorrJVF to build a two-dimensional likelihood ratio to further separate hard-scatter jets from pile-up jets. The variable  $R_{p_T}$  is defined as the scalar  $p_T$  sum of the tracks divided by the fully calibrated jet  $p_T$ , the tracks in this definition originate from the hard-scatter vertex. The variable CorrJVF is the corrected Jet Vertex Fraction (JVF). It is defined as a fraction of the  $\sum p_T$  of the tracks and corrected for the  $N_{\text{vertex}}$  dependent scalar  $\sum p_T$  of tracks associated with pile-up. The number of pile-up tracks per event is computed from all tracks associated with vertices other than the hard-scatter vertex.

## 8.4 *b*-tagging

A process very similar to the  $H^* \rightarrow WW$  signal is the production and decay of pairs of top quarks,  $t\bar{t}$ . These events are characterized by the presence of two  $b$ -jets along with two charged vector bosons, since the top quark decays instantaneously and almost exclusively decays into  $Wb$ . This results in a signature very similar to the ggF +jets or VBF signal process, where the Higgs subsequently decays into  $WW$ . Since the additional jets in  $t\bar{t}$  events originate from a  $b$ -quark, a  $b$ -tagging algorithm is applied to identify jets originating from  $b$ -quarks and filter out this background.

The ATLAS  $b$ -tagging algorithm [192, 193] aims to identify jets that originate from a  $b$ -quark, illustrated in Fig. 8.9. The mechanism enables the tagging of jets and their classification as  $b$ -jets through the hadronisation of  $b$ -quarks into  $b$ -hadrons, which have a relatively long lifetime and thus travel away from the interaction point before decaying further. The  $b$ -hadrons have a lifetime of around 1.5 ps and are able to travel a distance of around  $\langle c\tau \rangle \approx 450 \mu\text{m}$  before decaying. This  $b$ -hadron then decays further into charged particles and a secondary vertex is created with tracks originating from its decay products.



**Figure 8.9:** Illustration of the topology of an event at the primary vertex with a  $b$ -jet, showing typical features that are helpful in  $b$ -tagging: a secondary vertex reconstructed from the decay of a  $b$ -hadron, and the large impact parameters  $d_0$  of the tracks corresponding to the  $b$ -hadron decay products. Adapted from [12].

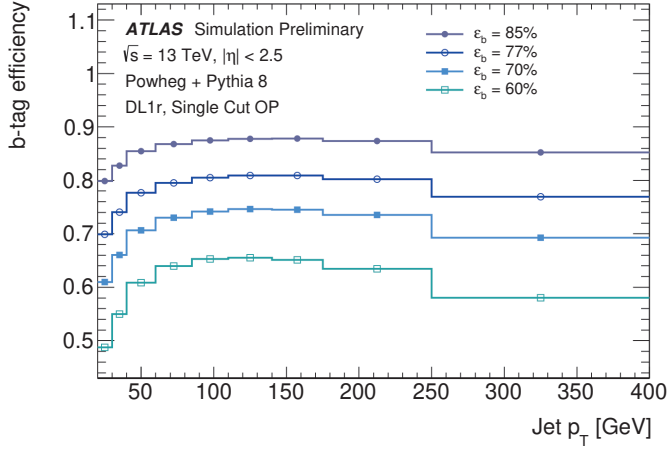
The tracks used for  $b$ -tagging are required to have  $p_T > 500$  MeV and have additional selection criteria applied, to reject potential fake tracks. The tracks are matched with hadronic jets satisfying  $p_T > 20$  GeV and  $|\eta| \leq 2.5$ , using the angular separation  $\Delta R$  between the track momenta and the jet axis. The  $\Delta R$  requirement is applied as a function of  $p_T$  to account for the boosted origin of the  $b$ -hadrons and varies between  $\Delta R = 0.45$  and  $\Delta R = 0.26$ , for jets with a  $p_T$  of 20 GeV, to 150 GeV, respectively. Along with this information, the location of the primary vertex is used to calculate the impact parameter of the tracks. Using this, displaced particle tracks are selected by requiring large transverse and longitudinal impact parameters. At last, jets are labelled as  $b$ -jets if they are matched to at least one weakly decaying  $b$ -hadron candidate with  $p_T > 5$  GeV. If no  $b$ -hadron candidate is found,  $c$ - and  $\tau$ -lepton candidates are searched for with different algorithms. If these are also not found, the jet is labelled as a light-flavour jet.

This last step is performed in a two-tiered approach, where first low-level algorithms attempt to reconstruct features of the  $b$ -hadron decay. The first low-level algorithm uses the impact parameters to define a probability for a track to originate from a  $b$ -hadron,  $c$ -hadron or light flavour jet and builds a discriminant observable from the probability.

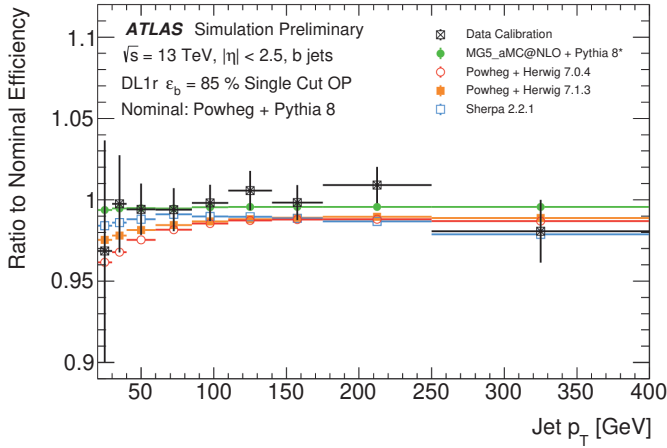
The next algorithm constructs the secondary vertex to identify all possible two-track vertices and then finds a single common displaced vertex associated with the heavy flavour decay, while rejecting tracks from long-lived particles such as a  $K_s^+(u\bar{s})$  or  $\Lambda^0(uds)$ . The last algorithm uses a Kalman filter to find the common line on which the primary and secondary vertex lie, by exploiting the topological structure of weak  $b$ - and  $c$ -hadron decays inside a jet.

Now the low-level variables are reconstructed, the information is fed to a multivariate classifier along with additional jet information. This algorithm is a neural network called DL1r [194], which is the default flavour tagging algorithm used in ATLAS Run-2 analysis. The neural network provides a multi-dimensional output which corresponds to the probability that the jet of interest is a  $b$ -jet,  $c$ -jet or light-flavour jet. The three probabilities are then combined in discriminants used to classify the jets. The efficiency of the  $b$ -tagging as a function of jet  $p_T$  is shown in Fig. 8.10. Several working points are defined, which correspond to the efficiency of the  $b$ -tagging algorithm.

The 85% efficiency working point, used throughout this thesis, provides the highest  $b$ -tagging efficiency. Calibration factors, shown in Fig. 8.11, correct for hadronization model differences across MC generators.



**Figure 8.10:**  $b$ -tagging efficiency for  $b$ -jets at the 85%, 77%, 70%, and 60% working points using the DL1r model [195].



**Figure 8.11:** Efficiency correction factors and data-to-MC scale factors for  $b$ -jets as a function of  $p_T$  for various Monte Carlo generators at the 85% working point. The error band on the data-to-MC scale factor includes systematic and statistical uncertainties. MC generator efficiency corrections show statistical uncertainties only [195].

## 8.5 Missing transverse momentum

Since the signature of interest in this thesis is the  $H \rightarrow WW^*$  decay, the final state particles are two leptons and two neutrinos. The neutrinos are weakly interacting particles and do not leave a signature in the detector. Therefore, the missing transverse momentum  $\vec{p}_T^{\text{miss}}$  is a crucial variable in this analysis.

Because the partons inside the colliding protons at the LHC collide parallel to the beam line, the total momentum in the transverse plane is zero. However, due to the presence of



neutrinos, finite resolution of the detector, the detector's limited acceptance or possible undetectable BSM particles, the total transverse momentum observed by the detector is not zero. Although, the direction of the carried away two momentum in the transverse plane is unknown the magnitude of the momentum can be calculated, because of momentum conservation, as the missing transverse momentum  $\vec{p}_T^{\text{miss}}$ . The  $\vec{p}_T^{\text{miss}}$  is calculated as the negative vector sum of the total transverse momentum in an event

$$\vec{p}_T^{\text{miss}} = - \sum_i \vec{p}_{T,i}.$$

As the mass of the outgoing undetected particles is unknown, it is often assumed to be zero. Using this assumption we can define  $\vec{p}_T^{\text{miss}}$  as the missing transverse energy defined as

$$E_T^{\text{miss}} = |\vec{p}_T^{\text{miss}}| = \left| - \sum_i \vec{E}_{T,i} \right|.$$

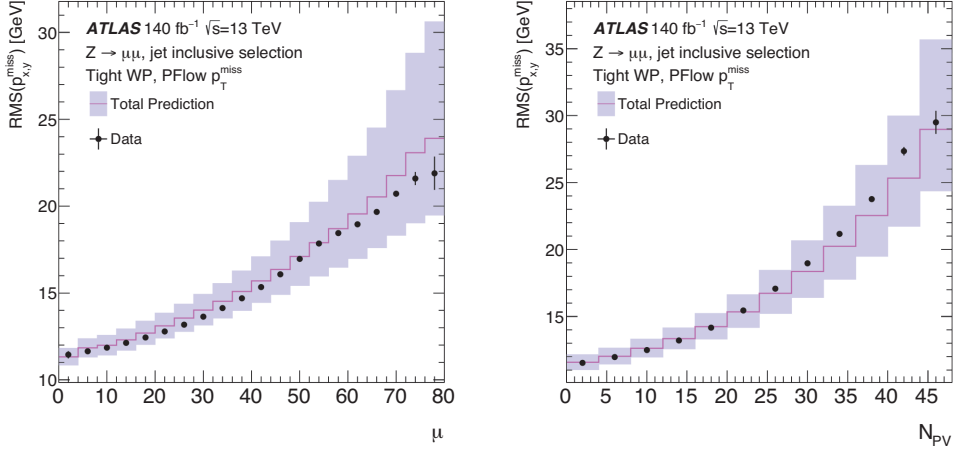
The missing transverse momentum reconstruction within ATLAS [196] consists of two main components. The first component takes into account all the hard event signals in the form of reconstructed and fully calibrated hard objects: these are electrons, muons,  $\tau$ -leptons, jets and photons. The second component is the soft event signals, which is all the activity in the detector not associated with the defined hard objects, but it is associated with the primary vertex: this comprises the soft jets, tracks and energy deposits in the calorimeters. The reconstruction of missing transverse momentum builds a set of observables from the 2-dimensional transverse momentum vectors  $\vec{p}_T = (p_x, p_y)$  of the various event objects given by

$$\vec{p}_T^{\text{miss}} = - \left( \underbrace{\sum_{\text{selected electrons}} \vec{p}_T^e + \sum_{\text{selected muons}} \vec{p}_T^\mu + \sum_{\text{selected } \tau\text{-leptons}} \vec{p}_T^\tau + \sum_{\text{selected photons}} \vec{p}_T^\gamma + \sum_{\text{selected jets}} \vec{p}_T^j}_{\text{hard term}} + \underbrace{\sum_{\text{unselected tracks}} \vec{p}_T^{\text{track}}}_{\text{soft term}} \right).$$

Along with  $\vec{p}_T^{\text{miss}}$ , the scalar sum of all transverse momenta  $\sum p_T = |\vec{p}_T|$  is also calculated as the overall scale for evaluating the hardness of the event, which is a measure of the overall activity within an event. In addition to these quantities, the azimuthal angle  $\phi^{\text{miss}} = \tan^{-1}(p_y^{\text{miss}}/p_x^{\text{miss}})$  of the missing momentum is calculated.

Before the  $\vec{p}_T^{\text{miss}}$  is calculated, an overlap removal procedure is performed which determines objects to add to the  $\vec{p}_T^{\text{miss}}$  calculation in cases where hard objects share one or more detector signal with each other. In this thesis, the track-soft-term (TST)  $E_T^{\text{miss}}$  reconstruction method is used which uses both track and calorimeter information in the hard term; the soft term only uses ID track information, but no calorimeter information; for robustness against pile-up and a better resolution, it does not include neutral objects or objects with  $|\eta| > 2.5$ . The *Tight* working point is used for the  $E_T^{\text{miss}}$  reconstruction in this thesis, which includes the usage of jets in the  $\vec{p}_T^{\text{miss}}$  calculation with  $p_T > 20$  GeV





**Figure 8.12:** The  $p_x^{\text{miss}}$  and  $p_y^{\text{miss}}$  resolution as a function of the pile-up  $\mu$  (left) and the number of primary vertices in the event. The resolution is shown for the TST  $E_T^{\text{miss}}$  reconstruction method using the *Tight* working point. A  $Z \rightarrow \mu\mu$  selection is applied to data and compared to  $Z \rightarrow \mu\mu$  MC events generated using Sherpa. The uncertainty band includes statistical, luminosity and detector uncertainties [196].

( $p_T > 30$  GeV) for  $|\eta| < 2.4$  ( $2.4 < |\eta| < 4.5$ ) along with JVT hard-scatter jet tagging. The resolution of  $p_x^{\text{miss}}$  and  $p_y^{\text{miss}}$  is shown in Fig. 8.12.

As there are multiple sources creating  $\vec{p}_T^{\text{miss}}$ , other than real non-interacting particles, such as particles disappearing in the beam pipe, mis-identified objects and pile-up, the variable called the missing transverse energy significance ( $E_T^{\text{miss}}$  significance) is defined. The  $E_T^{\text{miss}}$  significance, of which there are two types, is a measure of the compatibility of the  $E_T^{\text{miss}}$  with the hard objects in the event, instead of the mis-reconstruction of soft objects.

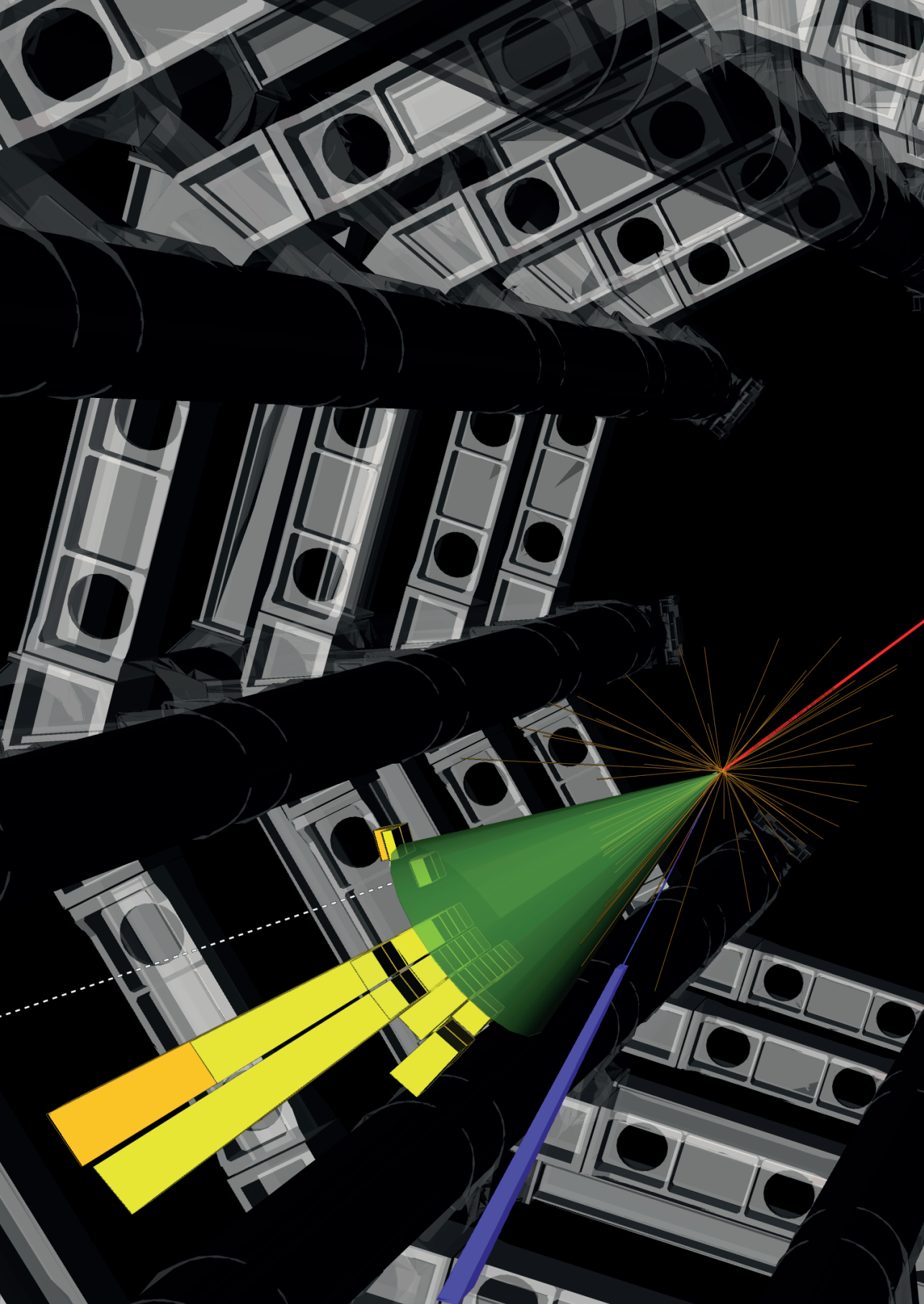
The object-based significance encodes the resolutions of the hard objects and correlations between them, combined in a likelihood function which encodes the possible mis-measurement of the objects used in the  $E_T^{\text{miss}}$  calculation. This likelihood model compares, per event, the hypothesis that no momentum is carried away by invisible particles ( $\vec{p}_T^{\text{miss}} = 0$ ) to the case of genuine momentum carried away by invisible particles ( $\vec{p}_T^{\text{miss}} \neq 0$ ). Another measure of this hypothesis is the event-based significance, which uses the scalar sum of the  $p_T$  of all hard objects in the event  $H_T$  and is defined as

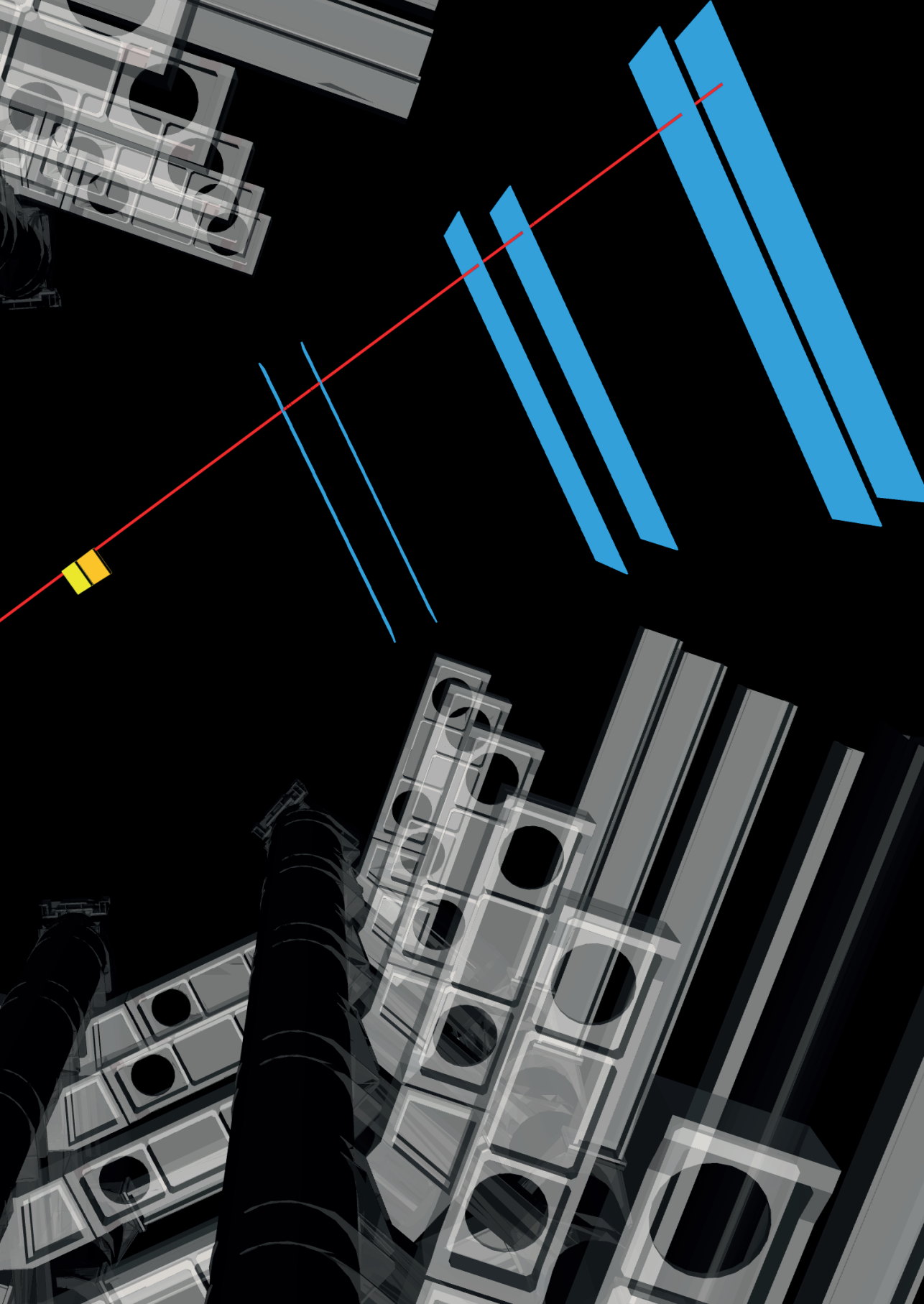
$$\text{event-based } \vec{p}_T^{\text{miss}} \text{ significance} = \frac{E_T^{\text{miss}}}{\sqrt{H_T}}.$$

Higher significance values indicate a better compatibility of the  $E_T^{\text{miss}}$  with the real  $\vec{p}_T^{\text{miss}}$  in the event. The  $E_T^{\text{miss}}$  significance is used in this thesis to suppress the background from mis-identified objects,  $Z \rightarrow ee/\mu\mu$  events and pile-up.

## PART III

### THE OFF-SHELL $H^* \rightarrow WW$ MEASUREMENT







## SIGNAL TOPOLOGY AND BACKGROUND ESTIMATION

---

This chapter presents the signal topology of the off-shell Higgs boson production in the  $H \rightarrow WW \rightarrow \ell\nu\ell\nu$  decay channel and the estimation of its most numerous backgrounds. After a brief motivation, it introduces the off-shell Higgs boson signal topology and background estimation techniques used in the off-shell analysis.

### 9.1 Motivation

The measured  $m_{VV}$  lineshape in an off-shell analysis is sensitive to anomalous HVV couplings and provides independent information on SM-EFT operators that affect the off-shell Higgs signal differently from its on-shell counterpart. This implies that measuring the  $\mu_{\text{off-shell}}$  signal strength is a promising probe for new physics. Moreover, the  $\mu_{\text{off-shell}}$  signal strength is directly proportional to the Higgs coupling modifiers without any dependence on the Higgs width, enabling broader interpretations. In contrast,  $\mu_{\text{on-shell}}$  directly depends on assumptions of the Higgs width.

Examples of analyses differentially probing the  $m_{VV}$  lineshape in the  $H \rightarrow WW \rightarrow \ell\nu\ell\nu$  decay channel using the Run-2 dataset are already present in the differential measurement of inclusive  $W^+W^-$  and  $W^+W^- + \geq 1$  jet production [197, 198]. The differential analyses provide excellent sensitivity to the properties of electroweak-boson self-interactions via anomalous triple-gauge-boson couplings and provide tests of perturbative QCD modelling of  $W^+W^-$  production. However, the effect of potential anomalous HVV couplings which could arise through the mediation of a non-SM off-shell Higgs boson is not targeted in the optimization.

The off-shell measurement presented in this thesis is optimized to be sensitive to the off-shell Higgs boson signal and its destructive interference with the  $W^+W^-$  continuum through higher-order (loop) corrections couplings to the high-mass  $W^+W^-$  system. It is therefore more optimized to probe new physics that may affect these processes.

In addition to providing direct Higgs coupling measurements, unique to off-shell Higgs measurements, the measurement is used to indirectly extract the value of  $\Gamma_H$ . This is achieved through a combined measurement of on-shell and off-shell Higgs boson production, as explained in § 3.4.3. This was successfully performed by ATLAS and CMS in recent years using the full Run-1 and Run-2 datasets [108, 199] and led to evidence for the presence of an off-shell Higgs using the  $H \rightarrow ZZ$  channel and constraining the total Higgs width.

The usage of the  $H \rightarrow WW$  channel to perform similar measurements is already motivated in 2014 [84] and has been used by the ATLAS experiment using Run-1 data in combination with the  $H \rightarrow ZZ$  channel [200]. Given the importance of the Higgs width in the SM, it is natural to investigate all channels used to perform an off-shell analysis to supply an orthogonal measurement. This orthogonal measurement can be used for increasing statistical precision via a combination.

On top of this, measuring the off-shell Higgs width in the  $H \rightarrow WW$  channel using charged ( $W^\pm$ ) vector bosons, rather than neutral ones ( $Z$ ), is compelling; as a discrepancy between the two width measurements could signal new physics. For example, in BSM models with additional scalars or a composite Higgs [201], the off-shell mechanisms for charged and neutral vector bosons can differ, leading to deviations from the SM expectation and the total Higgs width.

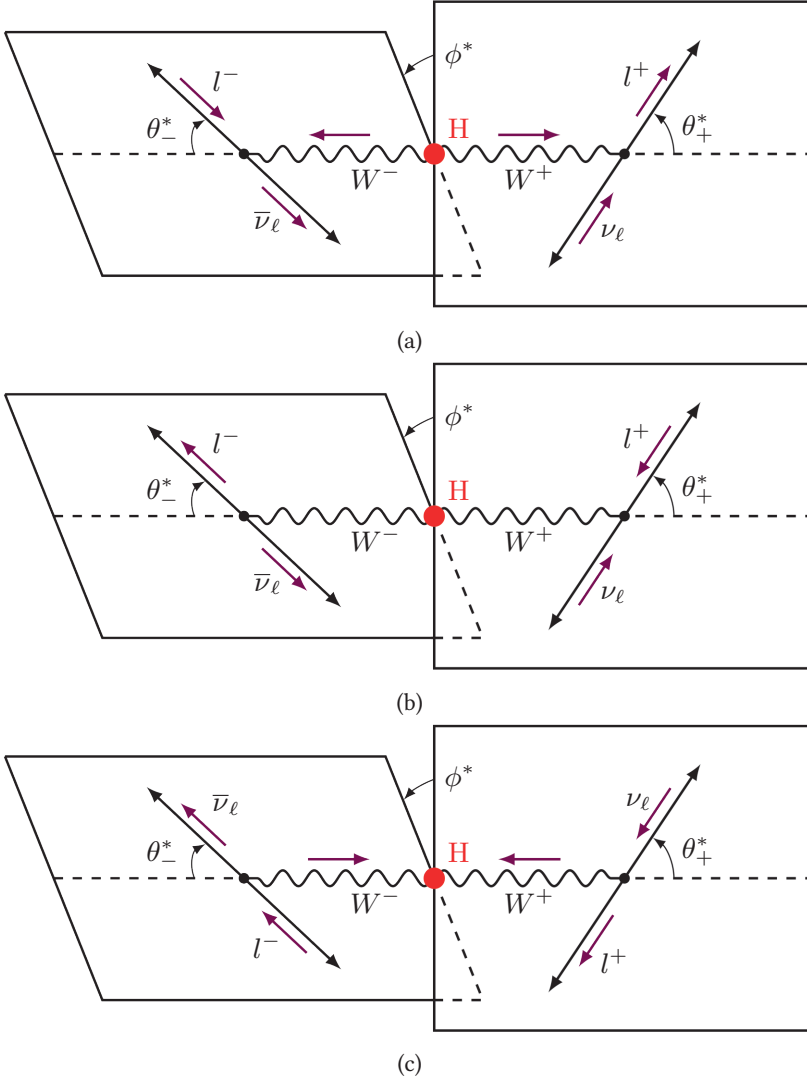
At first glance, a lack of mass resolution in the  $H \rightarrow WW \rightarrow \ell\nu\ell\nu$  channel appears to undermine the analysis method, which relies on a separation between the peak and off-shell regions. However, the crux of the method relies only on the existence of this separation and not on its exact reconstruction. Any variable that has well-defined on- and off-peak regions can be sensitive to the rescaling parameter  $\mu_{\text{off-shell}}$ , which is the POI in the off-shell analysis. Additionally, the  $WW$  channel also has advantages when compared to  $ZZ$ : the energy threshold of producing two real  $W$  bosons is lower compared to two real  $Z$  bosons, which allows for a larger phase-space to be probed and a larger number of events to be collected, both in terms of signal and interference with the  $W^+W^-$  continuum.

At last, a primary concern for the  $H^* \rightarrow WW$  measurement are the many backgrounds, which are numerous and more difficult to remove than in the  $ZZ$  channel. However, since the advent of neural networks within the field of particle physics, in combination with the continued improvement in background modelling techniques, the situation has changed. Nowadays, neural networks are able to reduce background contamination to a level where the signal can be extracted with greater accuracy. Compared to Run-1, background modelling also provides more reliable estimates with reduced uncertainties.

## 9.2 The off-shell Higgs boson signal topology

The off-shell Higgs boson signal topology in the  $H \rightarrow WW \rightarrow \ell\nu\ell\nu$  decay channel, depicted in Fig. 9.1, is characterized by the presence of two oppositely charged, isolated, high-energy leptons produced by two charged vector bosons.

This analysis considers both different-flavour ( $e\mu/\mu e$ ) and same-flavour ( $ee/\mu\mu$ ) channels of the  $H^* \rightarrow WW \rightarrow \ell\nu\ell\nu$  decay. Along with the two charged leptons, missing transverse energy is part of the topology, due to the neutrinos in the final state. In the SM, the Higgs boson is spinless, while  $W$  bosons carry spin 1. Hence, conservation of angular momentum allows three possible helicity configurations for  $W$  bosons in the Higgs rest frame.



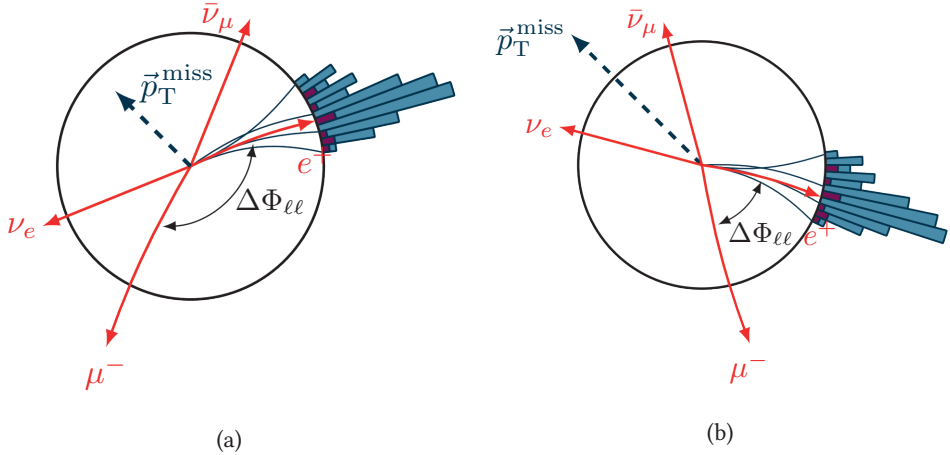
**Figure 9.1:** A schematic depiction of the  $H^* \rightarrow WW \rightarrow \ell\nu\ell\nu$  decay is shown, including the decay planes of the two leptonically decaying  $W$  bosons. Three different helicity configurations are illustrated from the perspective of the Higgs boson, which has no spin in the SM. These configurations are: (a) the positive  $W$  boson helicity configuration  $(+1, +1)$ , (b) the zero helicity configuration  $(0, 0)$ , and (c) the negative helicity configuration  $(-1, -1)$ . The decay planes of the two  $W$  bosons are depicted, showing the preferred orientation of the lepton and neutrino momenta. The angles  $\theta^*$  and  $\phi^*$  are defined as the angle between the direction of the leptons in the  $W$  boson rest frame and the direction of the  $W$  boson in the Higgs boson rest frame, respectively. Adapted from [12].



The three possible  $W$  boson helicity states ( $m_S = -1, 0, +1$ ) are: positive  $(+1, +1)$ , zero  $(0, 0)$ , and negative  $(-1, -1)$ . After the  $W$  bosons are produced, they decay leptonically via the V-A interaction, which is maximally parity-violating, into a lepton and a neutrino. Due to the V-A structure, the neutrinos have negative helicity and anti-neutrinos have positive helicity. This favours a lepton emission pattern where charged leptons align in the same direction while neutrinos recoil, creating substantial  $E_T^{\text{miss}}$ .

As a consequence of the spin and helicity structure of this decay topology, the opening angle between the charged leptons,  $\Delta\Phi_{ll}$ , tends to be small. Therefore, the invariant mass of the charged lepton pair,  $m_{ll}$ , is also small. More specifically, this occurs in the limit where  $\Delta\Phi_{ll} \rightarrow 0$  and  $m_{ll}$  approximates to  $\sqrt{E_{\ell^+} E_{\ell^-} (1 - \cos(\Delta\Phi_{ll}))} \rightarrow 0$ .

This effect is particularly significant in the on-shell region, where the invariant mass of the  $W$  bosons is required to be  $m_{WW} < 140 \text{ GeV}$ . As one moves into the off-shell region, where  $m_{WW} > 140 \text{ GeV}$ ,  $\Delta\Phi_{ll}$  tends to increase because of the larger invariant mass of the Higgs boson, which is depicted in Fig. 9.2.



**Figure 9.2:** A schematic depiction of the  $H^* \rightarrow WW \rightarrow e^+ \mu^- \nu_e \bar{\nu}_\mu$  in the  $\Phi$ -plane. The angle  $\Delta\Phi_{ll}$  is shown in (a) for a large opening angle and (b) for a small opening angle, resulting in a low missing energy and a high missing energy respectively. Adapted from [12].

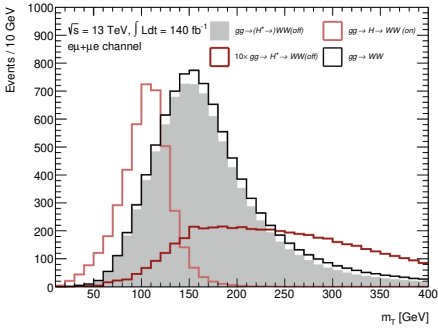
Unfortunately, the full kinematic reconstruction of the  $H \rightarrow WW \rightarrow \ell\nu\ell\nu$  decay is impossible due to the presence of two neutrinos in the final state. However, one can approximate the transverse mass  $m_T$  of the decaying Higgs system using

$$m_T = \sqrt{(E_T^{\ell\ell} + E_T^{\text{miss}})^2 - |\vec{p}_T^{\ell\ell} + \vec{E}_T^{\text{miss}}|^2}.$$

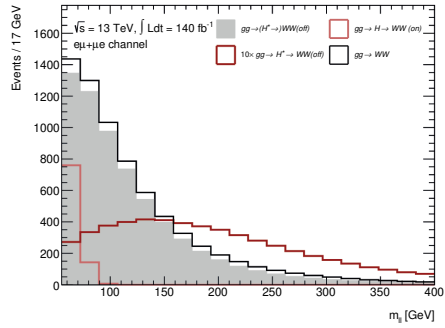
In this equation, the transverse energy of the di-lepton system is defined as  $E_T^{\ell\ell} = \sqrt{|\vec{p}_T^{\ell\ell}|^2 + m_{\ell\ell}^2}$ , where  $\vec{p}_T^{\ell\ell}$  represents the vector sum of the transverse momenta of the two charged leptons and  $\vec{E}_T^{\text{miss}}$  denotes the missing transverse energy. The reconstructed transverse mass provides a good approximation for the true transverse mass of the Higgs

boson system in the limit where the neutrinos are collinear with the charged leptons. In the analysis, it is used to distinguish the off-shell signal from various backgrounds, as it provides a natural cutoff around the mass of the on-shell boson, even though it is smeared by the resolution of the  $E_T^{\text{miss}}$ .

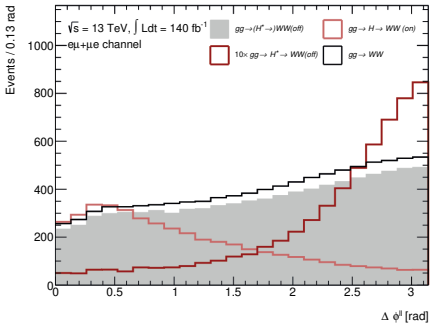
Using the aforementioned decay kinematics, the on-shell and off-shell Higgs processes are separated by one or more selections on the following observables:  $m_T$ ,  $m_{ll}$ ,  $\Delta\Phi_{ll}$ , and, to a lesser extent, the opening angle in the  $\eta$ -plane,  $\Delta\eta_{ll}$ . This separation is illustrated in the  $e\mu/\mu e$  channel in Figs. 9.3 to 9.6 for ggF and in Figs. 9.7 to 9.10 for VBF, after a loose preselection.



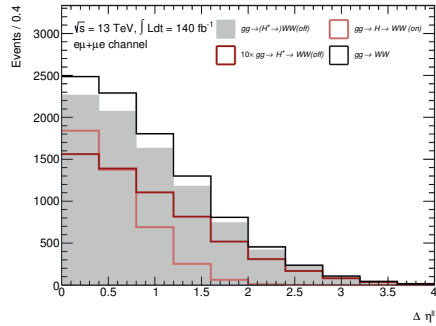
**Figure 9.3:** Distribution of the transverse mass,  $m_T$ .



**Figure 9.4:** Distribution of the invariant mass of the di-lepton system,  $m_{ll}$



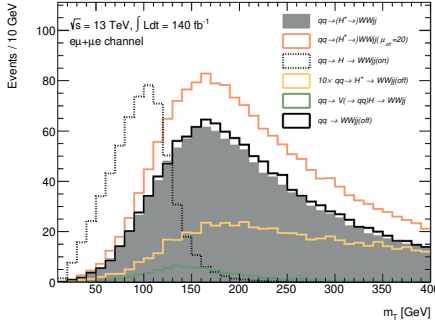
**Figure 9.5:** Distribution of the opening angle between the charged leptons in the  $\Phi$ -plane,  $\Delta\Phi_{ll}$ .



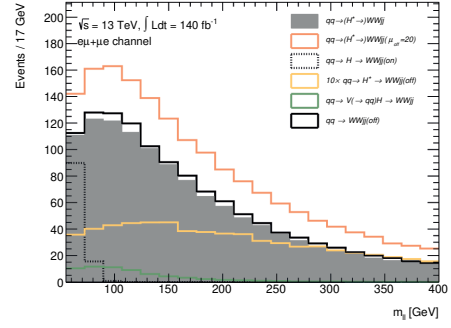
**Figure 9.6:** Distribution of the opening angle between the charged leptons in the  $\eta$ -plane,  $\Delta\eta_{ll}$ .

A selection of observables of the on-shell and off-shell ggF Monte Carlo samples in the different-flavour channel, after preselection, showing good separation power between on-shell and off-shell topologies.

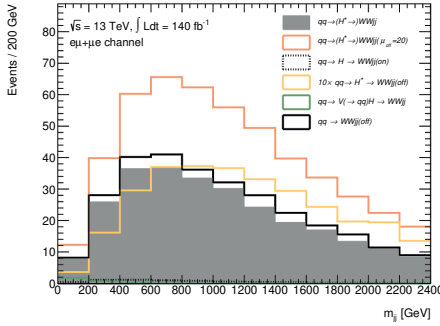
While the ggF analysis predominantly uses lepton kinematics to discriminate against major backgrounds, the VBF analysis benefits from additional topological features. In VBF production, the final state includes two jets originating from the radiation of the  $W$  bosons and the recoil against two outgoing quarks. This results in two energetic forward jets that arise from the hadronization of the outgoing quarks and are usually



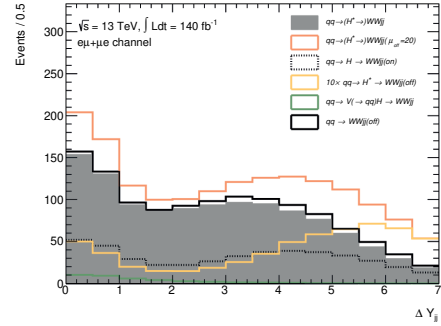
**Figure 9.7:** Distribution of the transverse mass,  $m_T$ .



**Figure 9.8:** Distribution of the invariant mass of the di-lepton system,  $m_l$ .



**Figure 9.9:** Distribution of the invariant mass of the two jets,  $m_{jj}$ .



**Figure 9.10:** Distribution of the rapidity gap between the jets,  $\Delta Y_{jj}$ .

A selection of observables of the on-shell and off-shell VBF fusion Monte Carlo samples in the different-flavour channel after preselection, showing good separation power between on-shell and off-shell topologies.

well separated in the  $\eta$ - $\phi$  plane.

As a result, the analysis can utilize the invariant mass of the two jets,  $m_{jj}$ , and the large rapidity separation between the two jets,  $\Delta\eta_{jj}$ , defined as  $|y_{j1} - y_{j2}|$ . Additionally, since there is no colour flow between the two jets and the Higgs boson is produced centrally, its decay products tend to lie between the two jets. All of these kinematic signatures are used in the VBF analysis to distinguish the signal from major backgrounds.

The Monte Carlo simulations considered as signal samples in the off-shell analysis are summarized in [Tabs. 9.1](#) and [9.2](#). These tables list the different production modes and decay channels of the generated signal samples. The simulation methods, generator settings and higher order corrections of these signal samples are extensively discussed in [§ 3.4](#).

**Table 9.1:** Summary of Monte Carlo samples used to simulate gluon-induced  $W^+W^-$  production. The simulated process, generator with PDF set, and  $K$ -factors are shown.

| Process                                                                          | Generator, PDF                 | $K(m_{WW}^{\text{truth}})$                                       |
|----------------------------------------------------------------------------------|--------------------------------|------------------------------------------------------------------|
| $gg \rightarrow \ell_{p,r}^+ \nu_{p,r} \ell_{p,r}^- \bar{\nu}_{p,r}, (p \neq r)$ |                                |                                                                  |
| $gg \rightarrow (H \rightarrow) WW \rightarrow \ell \nu \ell \nu$                | Sherpa+OpenLoops, NNPDF3.0nnlo | LO $\rightarrow$ NLO, 1.32 (NLO $\rightarrow$ N <sup>3</sup> LO) |
| $gg \rightarrow WW \rightarrow \ell \nu \ell \nu$                                |                                |                                                                  |
| $gg \rightarrow H \rightarrow WW \rightarrow \ell \nu \ell \nu$                  |                                |                                                                  |
| $gg \rightarrow \ell_{p,r}^+ \nu_{p,r} \ell_{p,r}^- \bar{\nu}_{p,r}, (p = r)$    |                                |                                                                  |
| $gg(\rightarrow H) \rightarrow WW/ZZ \rightarrow \ell \nu \ell \nu$              | Sherpa+OpenLoops, NNPDF3.0nnlo | LO $\rightarrow$ NLO, 1.32 (NLO $\rightarrow$ N <sup>3</sup> LO) |
| $gg \rightarrow WW/ZZ \rightarrow \ell \nu \ell \nu$                             |                                |                                                                  |
| $gg \rightarrow H \rightarrow WW/ZZ \rightarrow \ell \nu \ell \nu$               |                                |                                                                  |
| $gg \rightarrow (H \rightarrow) ZZ \rightarrow \ell \nu \ell \nu$                |                                | 1.32                                                             |

**Table 9.2:** Summary of Monte Carlo samples used to simulate quark-induced EW  $W^+W^- + 2$  jets production. The simulated process, generator and PDF set, cross-section, generator order,  $K$ -factor, and filter efficiency are listed.

| Process                                                                                                 | Generator, PDF                | Generator order |
|---------------------------------------------------------------------------------------------------------|-------------------------------|-----------------|
| $pp \rightarrow \ell_{p,r}^+ \nu_{p,r} \ell_{p,r}^- \bar{\nu}_{p,r} jj, (p \neq r)$                     |                               |                 |
| $pp \rightarrow (Hjj \rightarrow) WWjj \rightarrow \ell \nu \ell \nu jj$                                | MadGraph5+Pythia, NNPDF3.0nlo | 2j LO           |
| $pp \rightarrow (Hjj \rightarrow) WWjj \rightarrow \ell \nu \ell \nu jj, (\mu_{\text{off-shell}} = 20)$ |                               |                 |
| $pp \rightarrow WWjj \rightarrow \ell \nu \ell \nu jj$                                                  |                               |                 |
| $pp \rightarrow H \rightarrow WWjj \rightarrow \ell \nu \ell \nu jj$                                    |                               |                 |
| $pp \rightarrow \ell_{p,r}^+ \nu_{p,r} \ell_{p,r}^- \bar{\nu}_{p,r} jj, (p = r)$                        |                               |                 |
| $pp \rightarrow (Hjj \rightarrow) WWjj \rightarrow \ell \nu \ell \nu jj$                                | MadGraph5+Pythia, NNPDF3.0nlo | 2j LO           |
| $pp \rightarrow (Hjj \rightarrow) WWjj \rightarrow \ell \nu \ell \nu jj, (\mu_{\text{off-shell}} = 20)$ |                               |                 |
| $pp \rightarrow WWjj \rightarrow \ell \nu \ell \nu jj$                                                  |                               |                 |
| $pp \rightarrow H \rightarrow WWjj \rightarrow \ell \nu \ell \nu jj$                                    |                               |                 |

### 9.3 Backgrounds

The  $H^* \rightarrow WW$  channel is characterized by a large branching ratio, but is also affected by substantial backgrounds. Generally, two types of backgrounds are identified: reducible and irreducible backgrounds. Reducible backgrounds can be mitigated by applying a set of selection criteria; these backgrounds arise from processes that produce a slightly different final state as the signal and pass the selection criteria due to the misidentification of reconstructed objects. On the other hand, irreducible backgrounds stem from processes that have the same final state as the signal and pass the selection criteria due to the presence of genuine objects. In both the off-shell and on-shell analyses, the irreducible backgrounds are modelled using dedicated Monte Carlo simulations.

The approach to handling backgrounds in the off-shell and on-shell analyses is similar, with the primary distinction being the use of off-shell signal Monte Carlo samples, which ensures consistent background modelling across both analyses. For the most significant backgrounds, dedicated control and validation regions are established; these regions are

enriched with the target background process to ensure the accuracy of the modelling by comparing it with actual data. Reducible backgrounds are modelled using a data-driven fake factor technique; this method provides a reliable estimate of the reducible backgrounds that involve mis-identified leptons.

### 9.3.1 The $gg \rightarrow WW$ and $EW\ qq \rightarrow WWjj$ backgrounds

As detailed in § 3.4, the irreducible  $gg \rightarrow WW$  background strongly interferes with the off-shell ggF signal. It is therefore estimated along with the signal, using the S, B, and SBI samples, normalized to  $N^3\text{LO}$  in QCD with inclusive and differential K-factors. The sample only includes diagrams which are not covered by the  $qq \rightarrow WW$  background.

Similar to the  $gg \rightarrow WW$  process, the  $EW\ qq \rightarrow WWjj$  background strongly interferes with EW Higgs production and is therefore estimated alongside the signal using the S, SBI, and SBI20 samples at LO in QCD and EW.

### 9.3.2 The $qq \rightarrow WW$ background

The most important irreducible background for the ggF production mode is the  $qq \rightarrow WW$  background. It can be produced in association with jets, being prominent in the VBF backgrounds as well. It also interferes destructively with the VBF signal when produced in association with two jets.

The complete  $qq \rightarrow WW$  background used in the analysis is divided into processes containing four ( $\alpha_{EW}^4$ ) or six ( $\alpha_{EW}^6$ ) electroweak vertices and those containing a mixture of QCD and EW vertices, while also excluding diagrams with closed fermion loops, as not to overlap with gluon-induced  $WW$  production. They are generated using SHERPA 2.2.12 where NLO EW virtual corrections are applied to the nominal sample via the use of dedicated event weights [202]. The samples achieve NLO accuracy in QCD for up to one jet and LO for up to three jets. The  $qq \rightarrow WW$  normalization is estimated from data in the analysis.

### 9.3.3 The top quark background

Overall, the largest irreducible background where jets are required in the final state is when top quarks are involved. The top quark has a decay signature very similar to a  $W$  decay with an additional  $b$ -jet. The top quark decays, without hadronizing, into a  $b$ -jet and a  $W$  boson,  $t \rightarrow Wb$ , the  $W$  boson then decays leptonically, mimicking the signature of a single  $W$ -decay. The largest two processes involving top decays are the  $t\bar{t}$  and  $Wt$  processes. The  $t\bar{t}$  process is the production of a top-antitop pair, where both top quarks decay leptonically. The  $Wt$  process is the production of a single-top quark in association with a  $W$  boson, where the top quark decays leptonically. In order to reduce the contamination of top-related backgrounds, a veto is placed on the presence of  $b$ -quarks in the final state.

The  $t\bar{t}$  background is modelled using Powheg+Pythia at NLO in QCD, the A14 tune [113] for parton showering and hadronization, the NNPDF2.3lo PDF set [203], and normalized to data via dedicated control regions. In POWHEG-BOX v2, the  $h_{\text{damp}}$  parameter, which is a resummation damping factor that regulates high- $p_T$  radiation, is set to 1.5 times the mass of the top quark [204].

Previous studies show improved agreement between data and  $t\bar{t}$  predictions, especially for the top quark  $p_T$  distribution, when compared with NNLO calculations [205]. The  $t\bar{t}$  simulation is therefore further refined via a recursive reweighting procedure that adjusts the top quark  $p_T$ ,  $t\bar{t}$   $p_T$ , and  $t\bar{t}$ -mass distributions toward NNLO QCD precision [206].

The  $Wt$  and s-channel single-top processes are modelled using Powheg+Pythia at NLO in QCD with the A14 tune for parton showering, the NNPDF3.0nlo PDF set [207] and the five-flavour scheme. The single-top quark production in the t-channel is generated at NLO in QCD with Powheg+Pythia and the four-flavour scheme. All these top processes are normalized to data via the same control region as  $t\bar{t}$ . A diagram removal scheme [208] avoids overlap between  $tW$  and  $t\bar{t}$  samples, and event yields are scaled to an NLO QCD cross-section with NNLL soft-gluon corrections [209, 210].

#### 9.3.4 The $pp \rightarrow Z/\gamma^* + \text{jets}$ background

The Drell-Yan (DY) process  $Z/\gamma^* \rightarrow \ell\ell$ , produced in association with jets, is the most abundant background in the same-flavour channel and also features prominently in low mass regions in the different-flavour channel via the  $Z/\gamma^* \rightarrow \tau\tau \rightarrow \ell\nu\nu\ell\nu\nu$  decay. The DY process is modelled using Monte Carlo simulations at NLO in QCD and normalized to the data via the use of background-enriched regions. Additionally, NLO EW virtual corrections are applied via the inclusion of internal weights available in the latest version of SHERPA 2.2.11/2.2.14 [202]. The simulations, normalized to the NNLO QCD cross-section [211], are performed in the five-flavour scheme using Comix [212] and are NLO accurate for up to 2 jets and LO for up to 5 jets.

Because the DY process is challenging to model accurately, an additional validation region is used to check the modelling of the DY process. Along with the validation region, a two-point comparison with the SHERPA 2.2.1/2.2.2 prediction is made, which captures the different  $p_T$ -spectra of the simulations in an uncertainty. This is performed to ensure the DY process has sufficient flexibility in the fit and can adjust its shape and normalization in the high- $p_T$  regime, where the simulation is challenging and carries high resummation and parton shower uncertainties.

Electroweak  $Z$  production with two jets is generated using Herwig for  $Z \rightarrow e\bar{e}/\tau\tau$  and Herwig for  $Z \rightarrow \mu\mu$  [68, 213]. The simulation is accurate at NLO in QCD for up to two additional parton emissions and uses the MMHT2014 NLO [214] PDF set.

### 9.3.5 The other $VV(V)$ and $V/\gamma^*$ backgrounds

Now that the most important backgrounds are covered, the remaining backgrounds are the  $VV(V)$  backgrounds, where  $V$  can be any boson. These processes often pass selection through the misidentification of a  $\gamma$  for an electron through the creation of an  $e^-e^+$  pair. For the  $WZ, ZZ$  backgrounds there is the possibility to decay via two, three or four leptons plus the addition of neutrinos or quarks. If a lepton is lost in the reconstruction or via acceptance effects, these backgrounds may end up in the signal region. However, by applying a lepton veto of a possible third lepton, these backgrounds are suppressed.

The  $VV$  and  $V/\gamma^*$  backgrounds are modelled using SHERPA 2.2.12 and 2.2.8 respectively, with the exception of the semi-leptonic final states for  $ZZ$  and  $WZ$ , which were generated with SHERPA 2.2.11. The samples are accurate at NLO in QCD for up to one jet and LO in QCD for up to three jets.

Triboson ( $VVV$ ) processes are modelled with factorised gauge-boson decays in SHERPA 2.2.2, achieving NLO accuracy inclusively and LO for up to two additional jets. Loop-induced  $gg \rightarrow ZZ$  events are generated at LO in QCD with up to one extra parton using SHERPA 2.2.2.

### 9.3.6 Other Higgs production

Since the signal in the off-shell analysis consists of Higgs events with  $m_{WW} > 140$  GeV, the on-shell process is treated as a background. Other Higgs production modes ( $VH$ ,  $t\bar{t}H$ ,  $tH$ ) are also considered as backgrounds, though their yields are small. The most frequent Higgs decay in the signal regions, aside from  $H^* \rightarrow WW$ , is  $H \rightarrow \tau\tau$  via  $ggF$ , which shares a similar decay topology with  $H^* \rightarrow WW$ .

The  $H \rightarrow \tau\tau$  mode is simulated at NNLO QCD with POWHEG NNLOPS [215] and normalized to N<sup>3</sup>LO QCD and NLO EW cross sections [76]. The  $t\bar{t}H$  mode is also simulated with POWHEG-BOX v2 [215, 216] at NLO and is normalized to an NLO QCD and EW cross section [76].

### 9.3.7 The $pp \rightarrow W$ + jets and multi-jet backgrounds

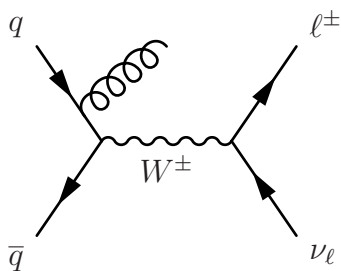
A large contribution of events that enter the signal region originate from  $W$  bosons produced in association with one or more jets, which mimic the signal topology due to the misidentification of jets as prompt leptons. A similar situation occurs when one or more purely hadronic jets is mis-identified as a charged lepton, often accompanied in combination with  $E_T^{\text{miss}}$ .

Other non-prompt leptons in an event originate from the leptonic decays of heavy quarks via a  $W$  boson, accompanied by a neutrino. Another possibility for creating fake leptons is the misidentification of a hadronic jet for an electromagnetic shower caused by a charged lepton. Since the mis-identified leptons are notoriously difficult to model, a

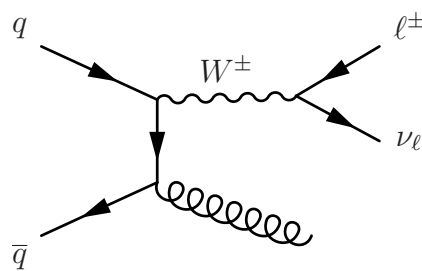
data-driven approach is employed to estimate this background in the analysis regions.

### 9.3.8 Mis-identified lepton estimation

Mis-identified leptons, also known as fake leptons, originate from the misidentification of jets as prompt leptons or from non-prompt leptons that mainly come from the decay of hadrons containing a heavy flavour quark. Two examples of  $W$ +jets Feynman diagrams are shown in Figs. 9.11 and 9.12. Due to the inaccurate modelling of fake leptons in Monte Carlo simulations, the data-driven fake factor method [217] is used to estimate the fake lepton background in both the off-shell and on-shell analyses.



**Figure 9.11:** The s-channel production of a  $W$  boson through quark fusion in association with a jet.



**Figure 9.12:** The t-channel production of a  $W$  boson through quark scattering in association with a jet.

To model misidentification using a Monte Carlo simulator, it must accurately predict all possible mechanisms. This is a challenging task, as the misidentification of jets as leptons is a complex process that depends on the detector response, the event topology and the kinematics of the jets. Furthermore, the identification algorithms used to reconstruct objects depend upon various parameters and are optimized for the best performance on data. In order to not depend on mis-modelling of either simulation, or reconstruction algorithms, the fake factor method is used to measure the number of events containing at least one object mis-identified as a lepton.

The approach is based on the assumption that the chance for an object to be mis-identified as a lepton depends on its inherent properties such as  $p_T$ ,  $\eta$ , the leptons isolation and identification criteria, defined in Tab. 9.3. The leptons which enter the signal region are required to pass a set of selection criteria, which are optimized to select prompt leptons called ID leptons (identified leptons). Other leptons, which cannot satisfy the ID criteria pass a looser set of selection criteria, called the anti-ID criteria (anti-ID leptons). The fake factor is applied to extrapolate the anti-ID yield in the anti-ID region to the ID yield in the ID-region.



**Table 9.3:** Requirements for ID and anti-ID leptons. Some channels require electrons to pass *LHMedium*, some require *LHTight*. Fake factors are provided for both options and the appropriate ones are used. The *AUTHOR* = 1 requirement ensures the electron candidate is not reconstructed as an ambiguous  $e/\gamma$  object, but exclusively as an electron.

| Electron                                           |                                  | Muon                                 |                              |
|----------------------------------------------------|----------------------------------|--------------------------------------|------------------------------|
| ID                                                 | anti-ID                          | UD                                   | anti-ID                      |
| $p_T > 15 \text{ GeV}$                             |                                  | $p_T > 15 \text{ GeV}$               |                              |
| $ \eta  < 2.47$ , excluding $1.37 <  \eta  < 1.52$ |                                  | $ \eta  < 2.5$                       |                              |
| $ z_0 \sin \theta  < 0.5 \text{ mm}$               |                                  | $ z_0 \sin \theta  < 0.5 \text{ mm}$ |                              |
| $ d_0 /\sigma(d_0) < 5$                            |                                  | $ d_0 /\sigma(d_0) < 3$              | $ d_0 /\sigma(d_0) < 15$     |
| Pass <i>LHTight/LHMedium</i>                       | Pass <i>LHLoose</i>              | Pass <i>Quality Tight</i>            | Pass <i>Quality Medium</i>   |
| Pass <i>PLIVTight</i> isolation                    |                                  | Pass <i>PLIVTight</i> isolation      |                              |
| <i>AUTHOR</i> = 1                                  |                                  |                                      |                              |
|                                                    | Veto against identified electron |                                      | Veto against identified muon |

A set of fake factors is derived by evaluating the following ratio for each kinematic bin of electrons and muons which satisfy the equation

$$N_{>0,\text{fake}}^{\text{ID,ID}} = F_2 \left( N^{\text{ID,anti-ID}} - N_{0,\text{fake}}^{\text{ID,anti-ID}} \right) + F_1 \left( N^{\text{anti-ID,ID}} - N_{0,\text{fake}}^{\text{anti-ID,ID}} \right) - F_1 F_2 \left( N^{\text{anti-ID,anti-ID}} - N_{0,\text{fake}}^{\text{anti-ID,anti-ID}} \right).$$

Each variable  $N$  refers to a number of events and  $F_1$  and  $F_2$  are the fake factors. The superscripts state which lepton candidate has the highest  $p_T$ , and if it passes the ID or anti-ID criteria. The subscript specifies whether only events without a fake lepton are considered. If no subscript is given, no additional selection based on truth information is made.

A key difference between the on-shell and the off-shell analysis is the trigger menu, which should be accounted for in the application of fake factors when the fake lepton is triggered. In general, one can distinguish two cases when the fake lepton is triggered: either only one lepton fired SLT and the other lepton did not fire the SLT, or both leptons fired the SLT. It is noteworthy in these occasions to state that the identification and isolation criteria of SLT are more stringent than that of anti-ID. This means that anti-ID leptons which fired SLT are less likely to be fakes than general anti-ID leptons, and applying nominal fake factors to such leptons causes a significant bias. Therefore, nominal fake factors need to be appropriately replaced by triggered fake factors when the SLT is fired.

Several of these cases require a dedicated fake factor derivation, visualized in [Fig. 9.13](#). These cases require to define two different fake factor sets, which is achieved by setting up the  $Z/\gamma^* + \text{jets}$  and di-jets fake factor measurement.

| $N_{l_1 l_2}$ | ID, ID        | AID, ID       | ID, AID       | AID, AID      |
|---------------|---------------|---------------|---------------|---------------|
| <b>RR</b>     | $N_{rr}^{ii}$ | $N_{rr}^{ai}$ | $N_{rr}^{ia}$ | $N_{rr}^{aa}$ |
| <b>FR</b>     | $N_{fr}^{ii}$ | $N_{fr}^{ai}$ | $N_{fr}^{ia}$ | $N_{fr}^{aa}$ |
| <b>RF</b>     | $N_{rf}^{ii}$ | $N_{rf}^{ai}$ | $N_{rf}^{ia}$ | $N_{rf}^{aa}$ |
| <b>FF</b>     | $N_{ff}^{ii}$ | $N_{ff}^{ai}$ | $N_{ff}^{ia}$ | $N_{ff}^{aa}$ |

Fake background  $N_{fakes \geq 1}^{ii} = N_{fr}^{ii} + N_{rf}^{ii} + N_{ff}^{ii}$

MC Estimation  $N_{MC}$

Single-fake Control Region

Double-fake Control Region

Canceled Terms

**Figure 9.13:** Illustration of the different types of fake lepton configurations in the off-shell analysis. **RR** denotes the case where both leptons are real, **RF** denotes the case where the first lepton is real and the second lepton is fake, **FR** and **FF** denote the cases where the first lepton is fake and the second real or fake, respectively.

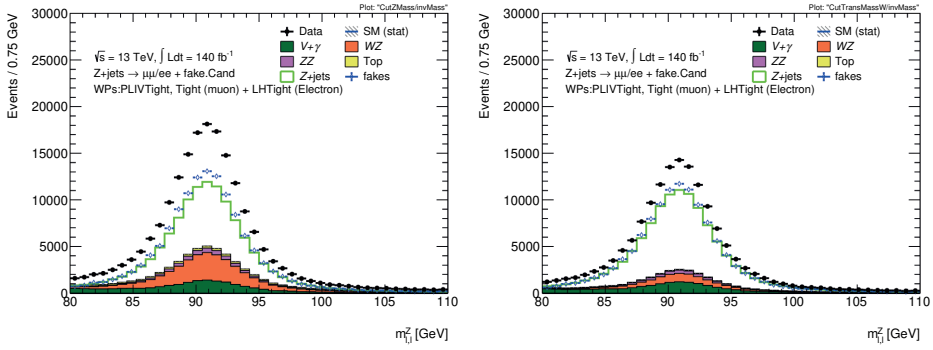
### $pp \rightarrow Z/\gamma^* + \text{jets}$ fake factor measurement

The nominal fake factors are calculated in regions enriched in fake leptons, by selecting events with at least one anti-ID lepton and requiring the presence of a second lepton. The fake factors are separately determined for electrons and muons, and are split in bins of  $p_T$  and  $\eta$  of the anti-ID lepton. For the estimation of non-triggered fakes,  $Z/\gamma^* + \text{jets}$  fake factors are used, in which one electron or muon pair with  $80 < m_{ll} < 110$  GeV or  $70 < m_{ll} < 110$  GeV is required, respectively. One of the leptons must fire a single-lepton trigger, such that the other leptons are unaffected by the trigger requirements. An extra requirement of  $M_T^W < 50$  GeV is applied to suppress contamination of prompt leptons in the fake enriched region.

The region with  $M_T^W > 50$  GeV is taken as a  $WZ$  control region in order to estimate the modelling of the most prominent prompt lepton background in the fake factor analysis. A normalization factor of  $1.01 \pm 0.01$  (*stat.*) is derived from the  $WZ$  control region and applied to the fake factor method, indicating sufficient modelling quality of the prompt leptons in this region. Distributions of the fake candidates for the  $Z/\gamma^* + \text{jets}$  fake factor measurement are shown in Fig. 9.14.

The yield of muon fake candidates with  $p_T > 50$  GeV is particularly small. In order to still arrive at the fake factor for this region, it is obtained by extrapolating from the  $25 < p_T < 50$  GeV bin via the use of Monte Carlo.

The parton flavour composition differs between  $W + \text{jets}$  and  $Z/\gamma^* + \text{jets}$  because of the weak boson coupling to quarks: for example,  $c$ -quarks are more commonly produced in



**Figure 9.14:** Distributions of the invariant mass of the reconstructed  $Z$  boson candidate at the stage of  $Z$  boson tagging (left) and after the  $m_T^W$  cut (right). The contributions without fake leptons are shown as filled histograms. The difference of the data and the filled histograms, which corresponds to the estimated fake lepton yield, is shown as blue points. The fake lepton yield predicted by the Powheg  $Z/\gamma^* + \text{jets}$  MC is shown in green for comparison.

association with  $W$  bosons via  $gs \rightarrow Wc$  couplings. With the estimation of a Monte Carlo-based flavour correction factor, the fake factors are corrected from the  $Z/\gamma^* + \text{jets}$  to the  $W + \text{jets}$  composition. The ratio of the fake factors is defined as the correction factor (CF) and calculated as:

$$CF = F_{MC}^Z / F_{MC}^W.$$

The value of the fake factors is then corrected by the correction factor, such that the fake factors in the analysis are used as extrapolated  $W + \text{jets}$  fake factors; this correction factor also comes with an uncertainty in the form of a 2-point Monte Carlo comparison.

At last, the uncertainty of the subtraction of the prompt lepton background in the fake enriched regions is estimated by evaluating the theoretical uncertainties on the prompt lepton simulation, which amounts to about 10%. This uncertainty is called the electroweak subtraction uncertainty. Using these distributions, an initial set of fake factors is derived by evaluating the following ratio for each kinematic bin of electrons and muons,

$$F = \frac{N^{\text{ID}} - N_{0, \text{fake}}^{\text{ID}}}{N^{\text{anti-ID}} - N_{0, \text{fake}}^{\text{anti-ID}}}. \quad (9.1)$$

The resulting fake factors of the  $Z \rightarrow \ell\ell + \text{jets}$  method, as well as impacts of the EW and sample composition uncertainties are shown in [Tab. 9.4](#).

**Table 9.4:** Fake factors estimated via the  $Z \rightarrow \ell\ell + \text{jets}$  method. Uncertainties are quoted as percentages of the central value. The statistical uncertainty includes the extrapolation uncertainty for the highest muon  $p_T$  bin. Total is the result of summing the individual uncertainties in quadrature. The EW subtraction and sample composition uncertainties are treated as correlated between different bins, while the statistical uncertainty is uncorrelated.

| Kinematic Region<br>ID and $p_T$ range | Nom.<br>Value | Stat. (%) | EW<br>Subtr. (%) | Sample<br>Comp. (%) | Tot. (%) |
|----------------------------------------|---------------|-----------|------------------|---------------------|----------|
| <b>Electron:</b>                       |               |           |                  |                     |          |
| LHMedium                               |               |           |                  |                     |          |
| 15.0 – 20.0 GeV                        | 0.041         | 7         | 5                | 9                   | 13       |
| 20.0 – 25.0 GeV                        | 0.048         | 13        | 10               | 15                  | 23       |
| 25.0 – 35.0 GeV                        | 0.074         | 13        | 10               | 24                  | 29       |
| 35.0 – $\infty$ GeV                    | 0.123         | 15        | 15               | 16                  | 27       |
| LHTight                                |               |           |                  |                     |          |
| 15.0 – 20.0 GeV                        | 0.027         | 9         | 8                | 16                  | 20       |
| 20.0 – 25.0 GeV                        | 0.033         | 17        | 12               | 31                  | 37       |
| 25.0 – 35.0 GeV                        | 0.054         | 17        | 12               | 20                  | 28       |
| 35.0 – $\infty$ GeV                    | 0.078         | 23        | 20               | 20                  | 36       |
| <b>Muon:</b>                           |               |           |                  |                     |          |
| Tight                                  |               |           |                  |                     |          |
| 15.0 – 20.0 GeV                        | 0.017         | 15        | 15               | 10                  | 23       |
| 20.0 – 25.0 GeV                        | 0.016         | 37        | 37               | 34                  | 62       |
| 25.0 – 50.0 GeV                        | 0.045         | 25        | 17               | 51                  | 60       |
| 50.0 – $\infty$ GeV                    | 0.074         | 28        | 17               | 51                  | 60       |

### Di-jets fake factor measurement

Fake factors obtained from a di-jets sample are used to estimate triggered fakes, i.e. fakes that cause an event to be triggered via a single-lepton trigger. Due to the high amount of fake events in di-jet events, the statistical uncertainties on these fake factors are considerably smaller. The di-jet events are triggered on an electron candidate by the prescaled HLT\_e12\_lhvloose\_nod0\_L1EM10VH trigger or on a muon candidate via the prescaled HLT\_mu14\_L1\_MU10 trigger. These triggers are available for the full Run-2 dataset and have a looser lepton selection such that they do not alter the ID/anti-ID ratio.

An event must satisfy the criteria in Tab. 9.5 to be included in the di-jets fakes estimation. Events with exactly one anti-ID or ID lepton object and at least one jet with  $p_T > 25$  GeV are considered as fake candidates. At least one of the jets must have an azimuthal separation with the lepton of  $\Delta\phi(\ell, \text{jet}) > 2.5$  in order to enforce the lepton+jet topology. The  $W$ +jets contribution, which has the largest cross-section among the processes with

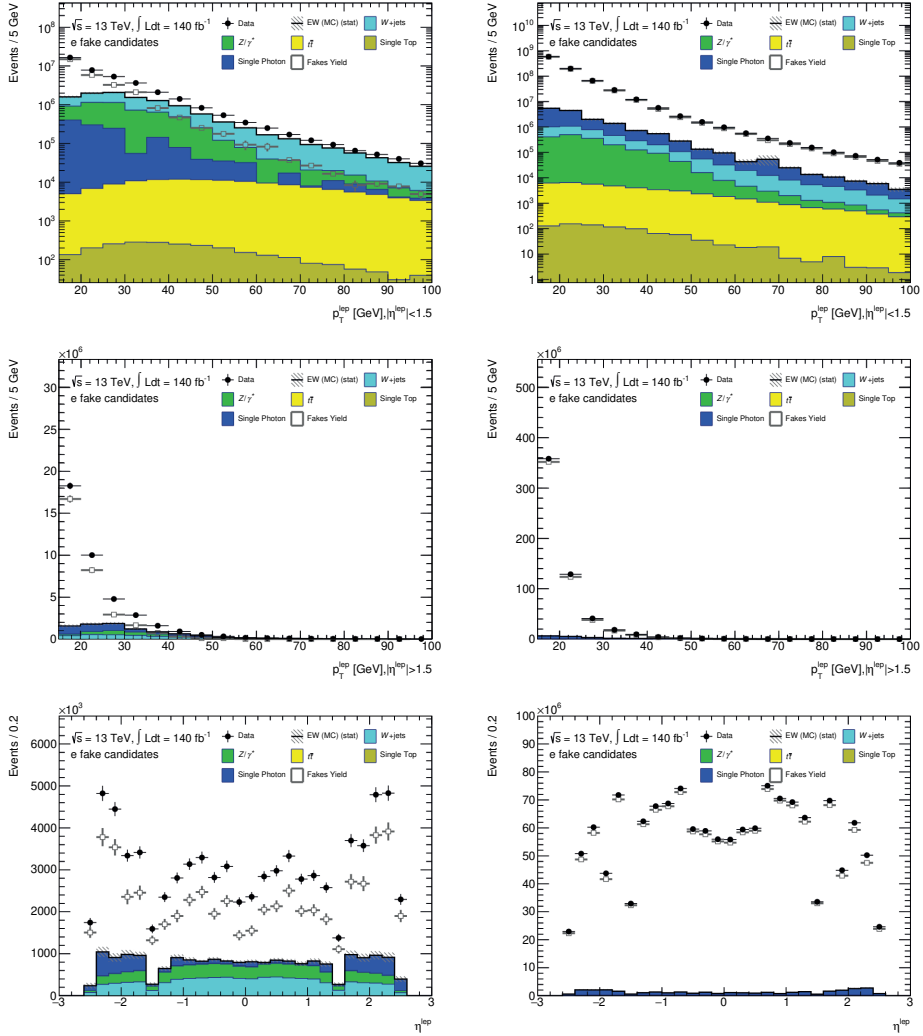
**Table 9.5:** Event selection used for the di-jets cutflow. Showing the ID and anti-ID regions, as well as the W+jet control region.

| ID/anti-ID region                               |                               | W+jet control region                       |
|-------------------------------------------------|-------------------------------|--------------------------------------------|
| ID                                              | anti-ID                       | ID                                         |
| pass HLT_mu14_L1_MU10 Trigger                   |                               | pass HLT_mu14_L1_MU10 Trigger              |
| pass HLT_e12_lhvloose_nod0_L1EM10VH Trigger     |                               |                                            |
| $\mu$ -trigger firing for $\mu$ -candidate      |                               | $\mu$ -trigger firing for $\mu$ -candidate |
| $e$ -trigger firing for $e$ -candidate          |                               | $e$ -candidates vetoed                     |
| <b>See Tab. 9.3<br/>for ID/anti-ID criteria</b> |                               | $ \eta  < 2.5$                             |
|                                                 |                               | $ z_0 \sin \theta  < 0.5 \text{ mm}$       |
|                                                 |                               | $ d_0 /\sigma(d_0) < 3$                    |
|                                                 |                               | Pass Quality Tight                         |
|                                                 |                               | Pass PLIVVeryTight isolation               |
| $p_T^{lep} > 15 \text{ GeV}$                    |                               | $p_T^{lep} > 120 \text{ GeV}$              |
|                                                 | $N_{jet} \geq 1$              |                                            |
|                                                 | $p_T^{jet} > 25 \text{ GeV}$  |                                            |
|                                                 | $\Delta\phi(\ell, jet) > 2.5$ |                                            |
|                                                 | $E_T^{miss} < 30 \text{ GeV}$ |                                            |
| $m_T^W < 30 \text{ GeV}$                        |                               | $30 \leq m_T^W < 50 \text{ GeV}$           |

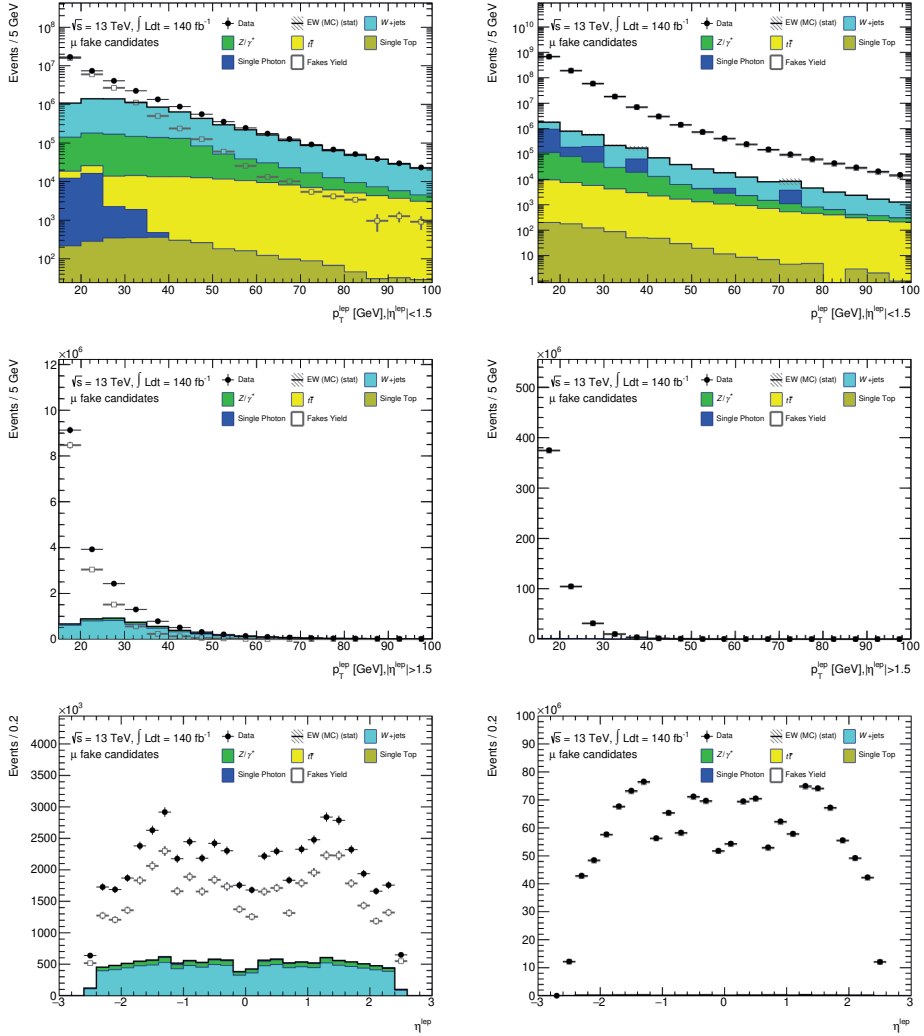
a prompt lepton, is suppressed via the cuts  $E_T^{miss} < 30 \text{ GeV}$  and  $m_T^W < 30 \text{ GeV}$  and its normalization estimated via a dedicated control region. Fakes in this control region are suppressed by requiring only muons satisfying the PLIV Very Tight isolation criteria. Additionally, a  $p_T^{lep} > 120 \text{ GeV}$  requirement is applied to suppress low- $p_T$  fakes. A sideband region spanning from  $30 < m_T^W < 50 \text{ GeV}$  is defined to extrapolate into the  $m_T^W < 30 \text{ GeV}$  ID and anti-ID regions.

Electron candidates are categorized as central ( $|\eta| < 1.5$ ) or forward ( $|\eta| > 1.5$ ), while muon candidates follow a slightly different definition as central ( $|\eta| < 1.05$ ) or forward ( $|\eta| > 1.05$ ). A  $W$ +jets normalization of  $0.902 \pm 0.016$  (*stat.*) and  $0.97 \pm 0.016$  (*stat.*) is extracted from the  $W$ +jets control region for the central and forward regions, respectively. These are then applied to the ID and anti-ID regions to get a better data to MC comparison for the fake estimation. The kinematic spectra of the fake candidates in the defined ID and anti-ID regions are shown in Fig. 9.15 and Fig. 9.16.

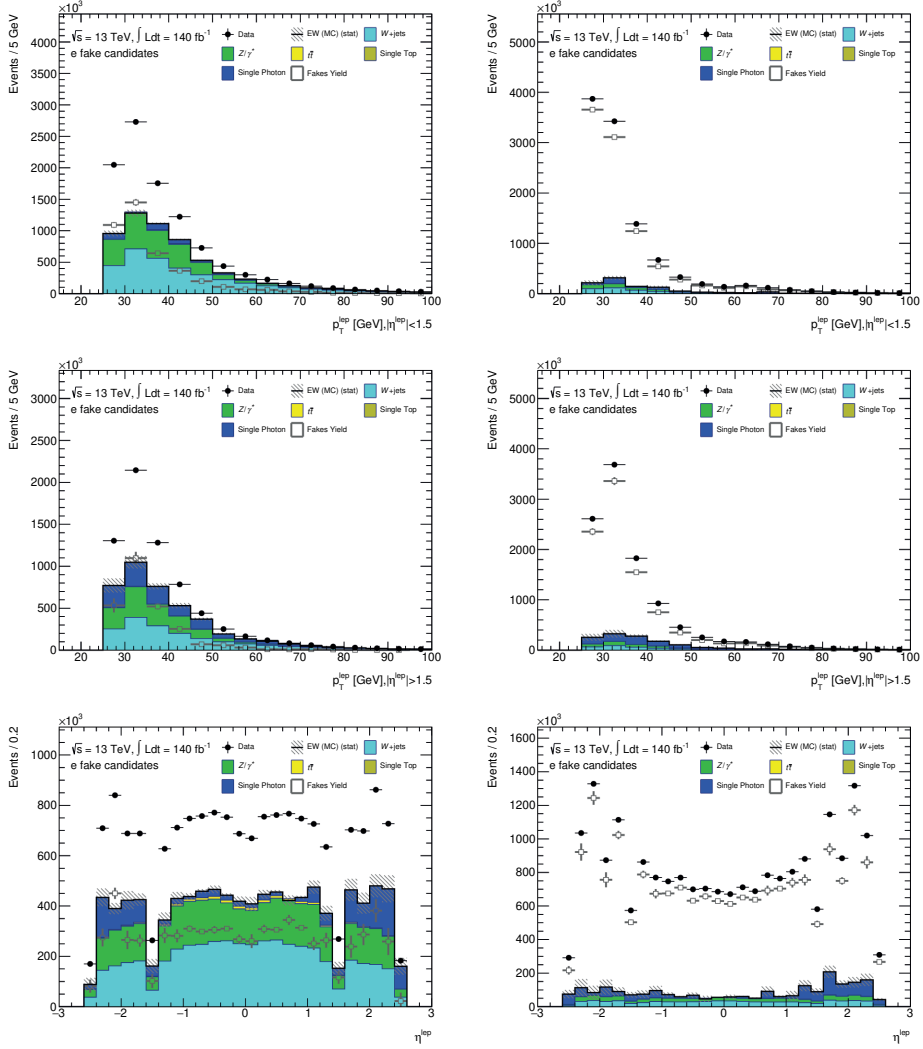
The fake factors for triggered fakes are evaluated with a stricter selection, which additionally requires the fake candidate to fire one of the single-lepton triggers in Tab. 6.2 and satisfy the corresponding  $p_T$  threshold. The ID and anti-ID regions after the trigger requirement are shown in Fig. 9.17 and Fig. 9.18 for electrons and muons, respectively. The fake leptons in the 2015 dataset are treated with an identical method but correspond to only  $< 3\%$  of the full dataset and are therefore not shown.



**Figure 9.15:** In the top row:  $p_T$  distributions for electron candidates with  $|\eta| < 1.5$  in the 2015-2018 dataset. The fake candidates pass the ID (anti-ID) criteria for the left (right) column. In the middle row the same distributions are shown for  $|\eta| > 1.5$ . The full  $\eta$  range is shown in the bottom row. The  $W$ +jets contribution is scaled using the respective normalization factor for the central and forward regions.

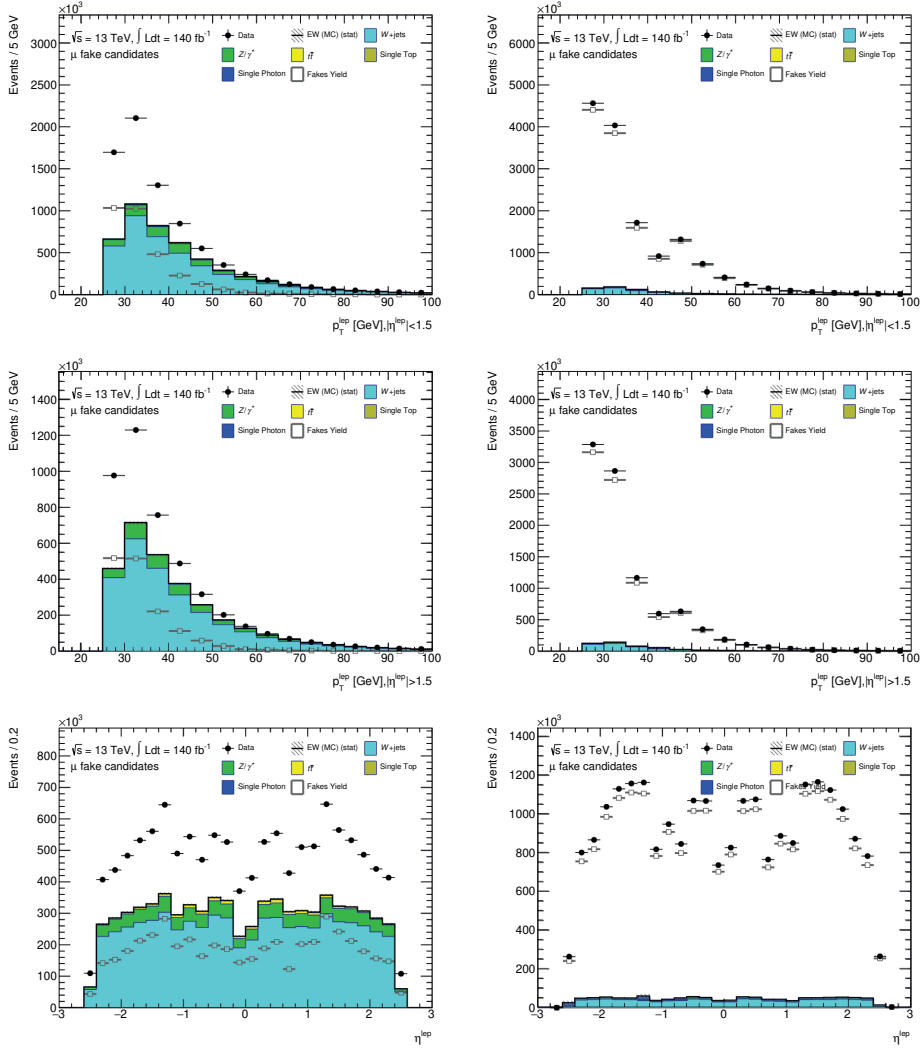


**Figure 9.16:** In the top row:  $p_T$  distributions for muon candidates with  $|\eta| < 1.05$  in the 2015-2018 dataset. The fake candidates pass the ID (anti-ID) criteria for the left (right) column. In the middle row the same distributions are shown for  $|\eta| > 1.05$ . The full  $\eta$  range is shown in the bottom row. The  $W$ +jets contribution is scaled using the respective normalization factor for the central and forward regions.



**Figure 9.17:** In the top row:  $p_T$  distributions for electron candidates with  $|\eta| < 1.5$  in the 2016-2018 dataset passing the triggers used in these years. The fake candidates pass the ID (anti-ID) criteria for the left (right) column. In the middle row the same distributions are shown for  $|\eta| > 1.5$ . The full  $\eta$  range is shown in the bottom row. The  $W$ +jets contribution is scaled using the respective normalization factor for the central and forward regions.





**Figure 9.18:** In the top row:  $p_T$  distributions for muon candidates with  $|\eta| < 1.05$  in the 2016-2018 dataset passing the triggers used in these years. The fake candidates pass the ID (anti-ID) criteria for the left (right) column. In the middle row the same distributions are shown for  $|\eta| > 1.05$ . The full  $\eta$  range is shown in the bottom row. The  $W$ +jets contribution is scaled using the respective normalization factor for the central and forward regions.

The triggered fake factors in the di-jets fake factor analysis are derived using the  $Z$ +jets fake factors and a trigger bias correction. The trigger bias correction, defined as  $R_{\text{trig}}$ , is derived from the ratio of the triggered and nominal fake factors. These are in turn defined as the ratio of the fake yields in the ID and anti-ID regions, where the superscript refers to the fake candidate.

$$R_{\text{trig}} = \frac{N_{\text{trig}}^{\text{ID}} - N_{0,\text{fake,trig}}^{\text{ID}}}{N_{\text{trig}}^{\text{anti-ID}} - N_{0,\text{fake,trig}}^{\text{anti-ID}}} \cdot \frac{N_{\text{nominal}}^{\text{anti-ID}} - N_{0,\text{fake,nominal}}^{\text{anti-ID}}}{N_{\text{nominal}}^{\text{ID}} - N_{0,\text{fake,nominal}}^{\text{ID}}}.$$

The trigger bias  $R_{\text{trig}}$  is less dependent on the sample composition than the absolute triggered fake factors themselves. Therefore, it has been decided to not directly use the fake factors for triggered fakes as derived from the di-jets analysis, for which at the time of this analysis no reliable flavour correction is possible. Instead, the nominal fake factors derived from the  $Z$ +jets analysis are multiplied with the  $R_{\text{trig}}$  derived in the di-jets analysis to obtain a fake factor for triggered fakes. Dedicated uncertainties are derived for  $R_{\text{trig}}$  and used for triggered fakes in addition to the uncertainties of the nominal fake factors. The fake factors for central triggered fake leptons are shown in Fig. 9.19 for the 2016-2018 dataset and summarized in Tab. 9.6. Identical tables are derived for the 2015 dataset and the forward region.

Since the sample composition uncertainty cannot be evaluated using di-jet Sherpa and Pythia8 Monte Carlo simulations due to limited statistics, a 40% uncertainty is added on top of the bias correction factors to account for extrapolation uncertainties from di-jets to  $W$ +jets. This uncertainty is estimated from the difference between the nominal  $Z$ +jets and di-jet fake factors.

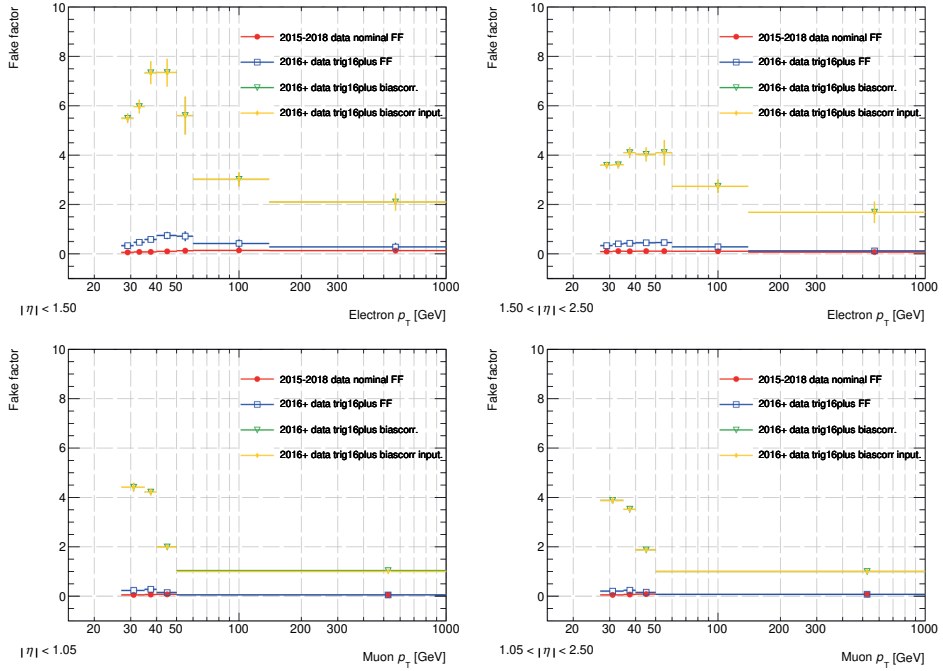
Furthermore, a 15% electroweak subtraction uncertainty is assigned to account for the extrapolation of the  $W$ +jets normalization factor derived in the dedicated control region. This uncertainty is increased to 25% for the low  $p_T$  regions below 60 GeV to account for the  $p_T$  extrapolation of the high  $p_T > 120$  GeV region into the ID and anti-ID regions. The relative size of the electroweak subtraction uncertainty on the final fake factor is small because of the high fakes purity of the ID and anti-ID region.

Fake triggered muons with  $p_T > 50$  GeV do not require a bias correction since no ID requirement is used on the triggers at this  $p_T$  threshold. Combining the aforementioned factors, bias corrections are computed separately for central and forward regions and applied to the electron and muon candidates independently.

### Application of fake factors

A unified fake factor application logic is developed, which is used by both the on-shell and off-shell analyses. This results in a common fake estimation, which facilitates a combination of the two channels.

If a channel uses single-lepton triggers (SLT) and di-lepton triggers (DLT), as in the on-shell analysis, then either the nominal fake factor or the nominal fake factor plus a



**Figure 9.19:** The Fake Factors and bias correction for electrons (top) and muons (bottom) for central (left) and forward (right) using the Tight ID Working point for electrons and Tight ID for muons. The uncertainties on the fake factors are the quadrature sum of the statistical and electroweak subtraction uncertainties. The triggered fake factors are derived using the 2016–2018 triggers in blue. The nominal fake factors are derived without a trigger requirement in red. The bias correction derived as the ratio between the triggered and nominal fake factors is shown in green. The bias correction input in the analysis is shown in yellow, where  $p_T$  bin for which we know there is not trigger bias are set to 1.

bias correction is applied, based on these rules: (1) For events selected by an SLT matched to an ID lepton, the nominal fake factor is used if the other lepton is anti-ID; (2) For any event passing a DLT, regardless of additional SLTs, the nominal fake factor and DLT bias correction for electrons is applied; (3) If selected by an SLT with an anti-ID lepton, the nominal fake factor and bias correction are applied to this lepton.

In channels using only SLT, as in the off-shell analysis, the rules are: (1) Events with only SLT matched to an ID lepton use the nominal fake factor if the other lepton is anti-ID; (2) For SLT-matched anti-ID leptons, the nominal fake factor plus bias correction is applied; if the other lepton is anti-ID, its nominal fake factor is used; (3) For both anti-ID leptons firing an SLT, the weight  $(F_1 \cdot F_2)(R_{\text{trig},1} \cdot R_{\text{trig},2} - R_{\text{trig},1} - R_{\text{trig},2})$  is applied, where  $F_1$  and  $F_2$  are fake factors, and  $R_{\text{trig},1}$  and  $R_{\text{trig},2}$  are bias corrections for leading and subleading-leptons, respectively, removing bias for the first lepton if the second fires the trigger and vice versa.

**Table 9.6:** Trigger bias corrections estimated via di-jets for the central region using the 2016-2018 dataset. The Statistical uncertainty includes the extrapolation uncertainty for the highest muon  $p_T$  bin. The EW subtraction uncertainty accounts for the subtraction of samples without fakes. Total is the result of summing the individual uncertainties in quadrature and depicted in percentage for illustrative purposes. The statistical analysis considers the uncertainties individually. The EW subtraction is treated as correlated between different bins, while the statistical uncertainty is uncorrelated.

| Kinematic Region<br>ID and $p_T$ range | Nom.<br>Value | Stat.   | EW<br>Subtr. | Tot.(%) |
|----------------------------------------|---------------|---------|--------------|---------|
| <b>Electron:</b>                       |               |         |              |         |
| LHMedium                               |               |         |              |         |
| 27.0 – 31.0 GeV                        | 4.39183       | 0.00235 | 0.00415      | 3.49    |
| 31.0 – 35.0 GeV                        | 4.58970       | 0.00395 | 0.00665      | 4.55    |
| 35.0 – 40.0 GeV                        | 5.60589       | 0.00582 | 0.01135      | 5.95    |
| 40.0 – 50.0 GeV                        | 5.73002       | 0.00885 | 0.02109      | 7.62    |
| 50.0 – 60.0 GeV                        | 4.49444       | 0.01771 | 0.02908      | 14.0    |
| 60.0 – 140.0 GeV                       | 2.74958       | 0.01455 | 0.05376      | 9.20    |
| 140.0 – $\infty$ GeV                   | 2.08608       | 0.01878 | 0.09706      | 15.7    |
| LHTight                                |               |         |              |         |
| 27.0 – 31.0 GeV                        | 5.49378       | 0.00197 | 0.00365      | 3.5     |
| 31.0 – 35.0 GeV                        | 5.96915       | 0.00328 | 0.00582      | 4.5     |
| 35.0 – 40.0 GeV                        | 7.33623       | 0.00491 | 0.00992      | 6.3     |
| 40.0 – 50.0 GeV                        | 7.34702       | 0.00753 | 0.01814      | 7.6     |
| 50.0 – 60.0 GeV                        | 5.60253       | 0.01544 | 0.02574      | 13.7    |
| 60.0 – 140.0 GeV                       | 3.02226       | 0.01274 | 0.04909      | 9.5     |
| 140.0 – $\infty$ GeV                   | 2.10160       | 0.01846 | 0.09529      | 16.6    |
| <b>Muon:</b>                           |               |         |              |         |
| Tight                                  |               |         |              |         |
| 27.0 – 35.0 GeV                        | 4.41525       | 0.00201 | 0.00560      | 3.9     |
| 35.0 – 40.0 GeV                        | 4.21963       | 0.00184 | 0.01570      | 3.4     |
| 40.0 – 50.0 GeV                        | 1.99632       | 0.00290 | 0.03024      | 4.6     |
| 50.0 – $\infty$ GeV                    | 1.            | 0.      | 0.           | 0.      |



## ANALYSIS OF OFF-SHELL HIGGS BOSON PRODUCTION USING $WW \rightarrow \ell\nu\ell\nu$ FINAL STATES

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This chapter presents the analysis of the off-shell Higgs boson production in the  $H^* \rightarrow WW \rightarrow \ell\nu\ell\nu$  decay channel [4]. After a brief introduction into the analysis philosophy, the trigger and object selection, the event selection and machine learning strategy are presented. This is followed by the uncertainty estimation, statistical model and results.

### 10.1 A brief overview of the analysis strategy

The event selection and categorisation in this analysis are based on the lepton flavour and the number of jets ( $N_{\text{jet}}$ ) in the event. The analysis is divided into two main categories: the  $ee/\mu\mu$  and  $e\mu/\mu e$  channels. The  $ee/\mu\mu$  channel is further divided into two categories: the electron-electron ( $ee$ ) and muon-muon ( $\mu\mu$ ) channels. The  $e\mu/\mu e$  channel follows a similar divide into the categories: electron-muon ( $e\mu$ ) and muon-electron ( $\mu e$ ). While the analysis framework processes the  $ee$ ,  $\mu\mu$ ,  $\mu e$ , and  $e\mu$  channels separately, the statistical interpretation merges  $ee$  with  $\mu\mu$  and  $\mu e$  with  $e\mu$ , in order to reduce the number of channels in the statistical model.

The analysis is performed in the following six regions, which all target the  $H^* \rightarrow WW$  decay mode in different configurations.

- $N_{\text{jet}} = 0$   $e\mu/\mu e$
- $N_{\text{jet}} = 1$   $e\mu/\mu e$
- $N_{\text{jet}} \geq 2$   $e\mu/\mu e$
- $N_{\text{jet}} = 0$   $ee/\mu\mu$
- $N_{\text{jet}} = 1$   $ee/\mu\mu$
- $N_{\text{jet}} \geq 2$   $ee/\mu\mu$

The analysis is split by flavour and jet multiplicity to account for different background compositions in each decay channel. This method allows for the optimization of the background rejection in each individual region and a better signal-to-background ratio in the overall analysis. Additionally, it allows to define data-driven normalization factors differentially in  $N_{\text{jet}}$ , in order to overcome potential mis-modelling in the simulation.

Most  $H^* \rightarrow WW$  events originate from the ggF production mode and primarily populate the  $N_{\text{jet}} = 0$  and  $N_{\text{jet}} = 1$  regions. When  $H^* \rightarrow WW$  events are produced via VBF, the production is accompanied by the radiation of two final state quarks, which form quark-initiated jets. Therefore, the VBF production tends to have higher jet multiplicities. However, because of the large rapidity gap between these jets, some may go missing

inside the forward part of the detector. This also results in a relatively high expected yield of ggF and VBF in the  $N_{\text{jet}} = 1$  region, which is accounted for in the optimization.

### 10.1.1 Triggers and object selection

Although the object candidate selection has been extensively discussed in [Chapter 8](#), a small summary of the objects used in this specific analysis is given here.

The charged leptons originating from the  $W$  boson are used as trigger signatures, with the full suite of triggers shown in [Tab. 6.2](#). The  $pp$ -collision data analysed in this measurement is taken with the ATLAS detector during the second run of the LHC from 2015-2018 at a centre-of-mass energy of  $\sqrt{s} = 13$  TeV. The ID requirements for electrons and muons are shown in [Tab. 9.3](#).

Jets containing  $b$ -hadrons are identified using the DL1r  $b$ -tagging algorithm. The  $b$ -jets are required to satisfy the requirements of the 85% efficiency working point determined with simulated  $t\bar{t}$  events. The jets originating from  $b$ -hadron decays with  $p_T \geq 20$  GeV and  $|\eta| < 2.5$  are referred to as  $b$ -jets in this analysis [[192](#), [193](#)].

An overlap removal procedure is applied to avoid double counting of objects, based on the angular separation  $\Delta R$  of the reconstructed objects. When two electrons have overlapping layer clusters in the calorimeters or a shared track, the lower- $p_T$  electron is removed. When a muon shares an ID track with an electron, the electron is removed. Also when a calorimeter-tagged muon shares an ID track with an electron, the muon is removed.

For electrons and jets, the jet is removed if it is within  $\Delta R < 0.2$  of the electron and the jet is not tagged as a  $b$ -jet. For muons and jets, the jet is removed if its within  $\Delta R < 0.2$  of the muon, the jet has less than three associated tracks with  $p_T > 500$  MeV, and the jet is not tagged as a  $b$ -jet. For any surviving jets, the lepton is removed if  $\Delta R(\text{jet}, \ell) < 0.4$ , with the goal of removing leptons from hadronic decays.

Jets are reconstructed using the anti- $k_t$  algorithm [[188](#)], with radius parameter  $R$  set to 0.4 implemented, in the FastJet package [[218](#)]. The inputs to the FastJet package are PFlow [[219](#)] objects in the central region ( $|\eta| < 2.5$ ) and topological clusters in the forward region ( $|\eta| > 2.5$ ). For jets with  $p_T > 25$  GeV and jet  $\eta < 4.5$ , a JVT *Tight* working point requirement is applied, corresponding to  $\text{JVT} > 0.5$  at  $20 < p_T < 60$  GeV and  $\eta < 2.4$ . A forward Jet Vertex Tagger (fJVT) *loose* working point requirement is also applied, corresponding to  $\text{fJVT} < 0.5$  at  $20 \text{ GeV} < p_T < 120 \text{ GeV}$  and  $2.5 < \eta < 4.5$ .

The  $E_T^{\text{miss}}$  is calculated as  $-\sum_i p_T^i$ : a negative sum of all the analysis objects. In addition, any leftover tracks which are not associated to already reconstructed objects are used in the analysis through the TST. The *Tight* working point of the  $E_T^{\text{miss}}$  is used [[220](#)].

## 10.2 Observable definitions

Several important common observables used in the  $H^* \rightarrow WW \rightarrow \ell\nu\ell\nu$  analyses for the separation of background and signal can be used simultaneously, in order to enhance the signal purity. These observables listed here are used in the event selection and in the multivariate analysis techniques. They are divided into three categories: background rejection, topological variables concerning the  $H^* \rightarrow WW \rightarrow \ell\nu\ell\nu$  decay, and VBF observables specific to this production mechanism.

### 10.2.1 Background rejection

The definition of the observables used to reduce the irreducible Drell-Yan and top backgrounds are:

- $m_{\tau\tau}$ : this variable is used to suppress the  $Z \rightarrow \tau\tau$  background, as well as the small contamination from  $H \rightarrow \tau\tau$ . Under the assumptions of the collinear approximation method [221], which assumes that the charged leptons seen are the products of the decay of a pair of  $\tau$  leptons, the neutrinos emitted in these decays are collinear with the charged leptons and these neutrinos are the only source of the observed  $E_T^{\text{miss}}$  in the event. This allows for the computation of the neutrino energy fractions and thus the invariant mass of the  $\tau$  pair.
- $N_{b\text{-jet}}$ : number of jets with  $p_T > 20$  GeV identified as  $b$ -jets from the  $b$ -tagged algorithm (DL1r) at 85% efficiency.
- $p_T^{\ell\ell}$ : transverse momentum of the di-lepton system. Related to the topological cuts but especially effective at rejecting Drell-Yan background.

### 10.2.2 Topological variables

Definitions of the topological variables used to enhance the Higgs signal and separate on-shell and off-shell signal are:

- $\Delta\phi_{\ell\ell}/\Delta\eta_{\ell\ell}$ : the products of the leptonic on-shell  $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$  decay tend to be more collimated than decay products from the non-resonant  $WW$  background because of the spin-zero initial state in the resonant process and the V-A structure of the weak interactions. The azimuthal angle  $\Delta\phi_{\ell\ell}$  and the pseudorapidity difference  $\Delta\eta_{\ell\ell}$  between the two leptons are used to exploit this feature.
- $m_{\ell\ell}$ : the invariant mass of the two leptons originating from the hard-scattering interaction. The off-shell signal tends to be at high  $m_{\ell\ell}$  for large  $m_H$ , for the same reason that  $\Delta\phi_{\ell\ell}$  is large for off-shell and small for on-shell.
- $m_T$ : the transverse mass is built from all the Higgs decay objects in the event.



### 10.2.3 VBF observables

Due to the specific kinematic signature of the VBF events, additional variables are available for VBF compared to the ggF signature, for performing event selection and as input for the multivariate analysis techniques. These variables involve the kinematics of the two highest- $p_T$  jets and are defined as:

- $\Delta y_{jj}$ : the signal is characterized by a separation of the two jets in rapidity, hence a gap in  $\Delta y_{jj}$  is a useful quantity to separate it from the backgrounds.
- Central-jet  $p_T$ : the central jet is defined as the jet with the largest  $p_T$  among all jets that lie within the  $\Delta y_{jj}$  of the two jets. The transverse momentum of the central jet (the central-jet  $p_T$ ) is used to define the central-jet veto (CJV):

$$\text{CJV} \begin{cases} = 1 & : \text{no central jet with } p_T < 30 \text{ GeV is present in the event,} \\ = 0 & : \text{a central jet with } p_T > 30 \text{ GeV is present in the event.} \end{cases}$$

- $m_{jj}$ : a high invariant mass of the jets ( $m_{jj}$ ) is a distinctive signature of the signal events.
- $\eta_{\text{lep}}$  centrality: this observable quantifies the exact positions of the leptons with respect to the two jets in the  $\eta$ -plane:

$$\begin{aligned} \text{OLV}_{l_0} &= 2 \cdot \left| \frac{\eta_{l_0} - \bar{\eta}}{\eta_{j_0} - \eta_{j_1}} \right| \\ \text{OLV}_{l_1} &= 2 \cdot \left| \frac{\eta_{l_1} - \bar{\eta}}{\eta_{j_0} - \eta_{j_1}} \right| \\ \eta_{\text{lep}} \text{ centrality} &= \text{OLV}_{l_0} + \text{OLV}_{l_1} \end{aligned} \quad (10.1)$$

where  $\bar{\eta} = (\eta_{j_0} + \eta_{j_1})/2$  is the average  $\eta$  of the two jets. The value of the outside-lepton-veto (OLV) observable is interpreted as:

$$(\text{OLV})_l \begin{cases} = 0 & : \text{the lepton is in the middle of the } \Delta y_{jj} \text{ between the two jets,} \\ < 1 & : \text{the lepton lies within the } \Delta y_{jj} \text{ between the two jets,} \\ > 1 & : \text{the lepton is outside the } \Delta y_{jj} \text{ between the two jets.} \end{cases}$$

### 10.3 Event selection

A set of selection criteria, based on the previously defined observables, is established to isolate  $H^* \rightarrow WW$  signal events while minimizing background contamination, then applied to both data and simulation. The selection criteria are grouped into four categories: *pre-selection*, *jet and flavour category*, *orthogonality requirements* and *background rejection*. The complete selection criteria for the  $e\mu/\mu e$  and  $ee/\mu\mu$  signal regions are shown in [Tab. 10.1](#).

**Table 10.1:** Event selection criteria for the  $e\mu/\mu e$  and  $ee/\mu\mu$  channels. The jet categorisation is based on the number of jets in the event. The orthogonality requirements are used to ensure that the signal regions are orthogonal to the control regions, the on-shell  $H \rightarrow WW^*$  measurement and the off-shell  $H^* \rightarrow ZZ$  measurement. The background rejection requirements are used to suppress the background.

| $e\mu/\mu e$ channel | 0-jet                                                                                                                                                                                                                                                              | 1-jet                                                             | $\geq 2$ -jet                   |
|----------------------|--------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------|-------------------------------------------------------------------|---------------------------------|
| Pre-selection        | Trigger selection and matching<br>two different-flavour oppositely-charged leptons<br>No additional leptons<br>$p_T^{\text{lead}} > 27 \text{ GeV}$ and $p_T^{\text{sublead}} > 15 \text{ GeV}$                                                                    |                                                                   |                                 |
| Jet categorisation   | $n_{\text{jets}} = 0$                                                                                                                                                                                                                                              | $n_{\text{jets}} = 1$                                             | $n_{\text{jets}} \geq 2$        |
| Orthogonality        | $m_{\ell\ell} > 100 \text{ GeV}$ or<br>$m_{\ell\ell} > 55 \text{ GeV}$ and $\Delta\phi(\ell, \ell) > 2.0$                                                                                                                                                          | $m_{\ell\ell} > 80 \text{ GeV}$<br>$\Delta\phi(\ell, \ell) > 1.8$ | $m_{\ell\ell} > 70 \text{ GeV}$ |
| Background rejection | $n_{b\text{-jets}} = 0$<br>$\Delta\eta(\ell, \ell) < 1.8$                                                                                                                                                                                                          |                                                                   |                                 |
|                      |                                                                                                                                                                                                                                                                    |                                                                   | CJV<br>OLV                      |
| $ee/\mu\mu$ channel  | 0-jet                                                                                                                                                                                                                                                              | 1-jet                                                             | $\geq 2$ -jet                   |
| Pre-selection        | Trigger selection and matching<br>two same-flavour oppositely-charged leptons<br>No additional leptons<br>$p_T^{\text{lead}} > 27 \text{ GeV}$ and $p_T^{\text{sublead}} > 15 \text{ GeV}$<br>$p_T^{\ell\ell} > 40 \text{ GeV}$<br>$\mathcal{S}_{\text{miss}} > 4$ |                                                                   |                                 |
| Jet categorisation   | $n_{\text{jets}} = 0$                                                                                                                                                                                                                                              | $n_{\text{jets}} = 1$                                             | $n_{\text{jets}} \geq 2$        |
| Orthogonality        | $\Delta R(\ell, \ell) > 1.8$<br>$m_{\ell\ell} > 55 \text{ GeV}$<br>$\Delta\phi(\ell, \ell) > 1.8$                                                                                                                                                                  |                                                                   | $m_{\ell\ell} > 70 \text{ GeV}$ |
| Background rejection | $n_{b\text{-jets}} = 0$<br>$ m_{\ell\ell} - 91 \text{ GeV}  > 15 \text{ GeV}$                                                                                                                                                                                      |                                                                   |                                 |
|                      |                                                                                                                                                                                                                                                                    |                                                                   | CJV<br>OLV                      |

In the pre-selection, the events are required to pass the trigger and trigger matching requirements. The  $p_T$  thresholds of the lowest triggers vary from year to year, shown in Tab. 6.2. The leading-lepton is required to have a  $p_T$  of at least 27 GeV, corresponding to the highest trigger threshold, and the subleading-lepton is required to have a  $p_T$  of at least 15 GeV in order to suppress Drell-Yan backgrounds. When a third charged lepton is present in the event, the event is rejected. When  $ee/\mu\mu$  leptons are identified, an additional  $p_T^{\ell\ell} > 40 \text{ GeV}$  requirement is applied in order to reject  $Z/\gamma^* + \text{jets}$  events, which are challenging to model due to resummation effects at low- $p_T$ . On top of this, a  $\mathcal{S}_{\text{miss}} > 4$  requirement is applied in order to suppress the  $Z/\gamma^* + \text{jets}$  background even further.

In the  $N_{\text{jet}} = 0$   $e\mu/\mu e$  channel, events satisfying  $m_{\ell\ell} > 100 \text{ GeV}$  or  $m_{\ell\ell} > 55 \text{ GeV}$  and  $\Delta\phi(\ell, \ell) > 2.0$  enter the signal regions. This selection targets the high invariant mass of

the off-shell  $WW$  system, while at the same time maintaining orthogonality with the on-shell analysis and its  $WW$  control regions. In the  $N_{\text{jet}} = 1$   $e\mu/\mu e$  channel, events are required to have  $m_{\ell\ell} > 80$  GeV and  $\Delta\phi(\ell, \ell) > 1.8$  to enter the signal regions, also targeting high invariant mass and large opening angles between the leptons. Events that enter the  $N_{\text{jet}} \geq 2$   $e\mu/\mu e$  channel, are required to have  $m_{\ell\ell} > 70$  GeV and pass the CJV and OLV cuts. On top of this, all regions feature a  $b$ -jet veto in order to suppress the  $t\bar{t}$  and  $Wt$  backgrounds. At last, the  $N_{\text{jet}} = 0$   $e\mu/\mu e$  channel has an additional  $\Delta\eta(\ell, \ell) < 1.8$  requirement in order to suppress the  $qq \rightarrow WW$  background.

In the  $N_{\text{jet}} = 0$   $ee/\mu\mu$  channel, after events have passed the preselection, a  $\Delta R(\ell, \ell) > 1.8$  selection is made to ensure orthogonality with the  $H^* \rightarrow ZZ$  off-shell measurement. This also suppresses  $Z/\gamma^* + \text{jets}$  by selecting events with a large opening angle in the  $\eta$ - $\Phi$ -plane. Furthermore, the off-shell signal is targeted by requiring  $m_{\ell\ell} > 55$  GeV and  $\Delta\phi(\ell, \ell) > 1.8$ .

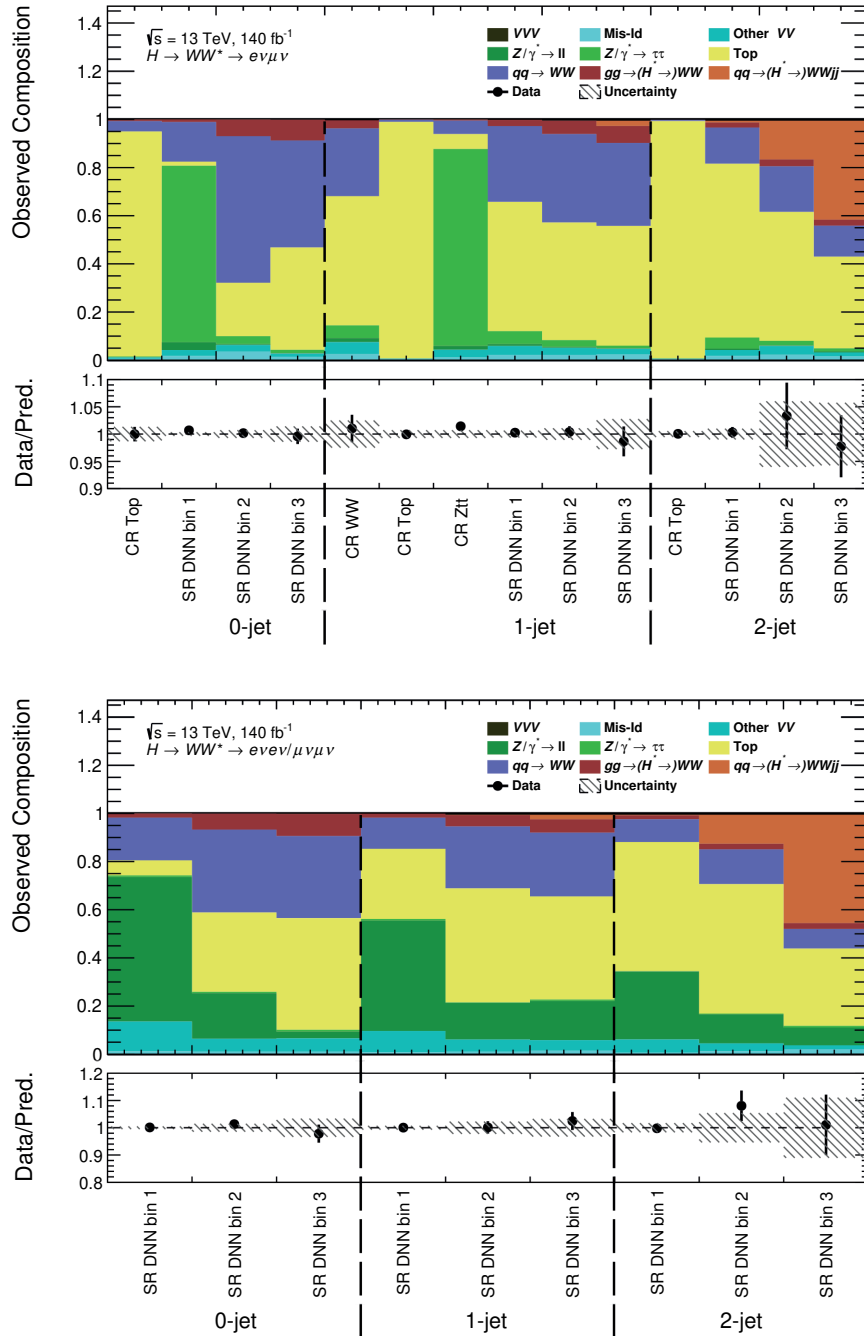
In the  $N_{\text{jet}} = 1$   $ee/\mu\mu$  channel, the  $m_{\ell\ell}$  requirement is increased to  $m_{\ell\ell} > 70$  GeV in order to ensure orthogonality with the  $ee/\mu\mu$  on-shell VBF analysis. After the orthogonality requirements have been passed, a  $b$ -veto is applied and events produced on the  $Z$ -peak are rejected by requiring  $|m_{\ell\ell} - 91\text{GeV}| > 15$  GeV. Similar to the  $N_{\text{jet}} \geq 2$   $e\mu/\mu e$  channel, the  $N_{\text{jet}} \geq 2$   $ee/\mu\mu$  channel features a CJV and OLV requirement.

### 10.3.1 Signal extraction

In order to extract the off-shell  $H^* \rightarrow WW$  signal yield, a binned profile likelihood fit is performed in a total of 23 regions simultaneously, with the inclusion of experimental and theoretical uncertainties as nuisance parameters. The 23 regions consist of 18 signal regions, varying in signal purity, and five control regions. A schematic overview of all signal and control regions is shown in Fig. 10.1.

In each signal region, the fit discriminant is chosen to be a proxy for the  $H^* \rightarrow WW$  invariant mass,  $m_{WW}^{\text{proxy}}$ , which shows high correlation with the real invariant mass,  $m_{WW}^{\text{particle}}$ , in the various signal regions.

In the control regions, a single-bin fit is employed in order to extract the normalization factors for the various backgrounds, without making any data-driven assumptions on the shape of the background. The 18 signal regions originate from the six regions of flavour and jet multiplicity. Each of these six regions features its own dedicated deep neural network, which is trained to separate the off-shell signal from the non-interfering background. The resulting probability score calculated by the deep neural networks is split into a *background-enriched* (DNN bin 1), *mixed* (DNN bin 2) and *signal-enriched* (DNN bin 3) region. All the regions are regarded as signal regions in the final fit.



**Figure 10.1:** A schematic overview of the  $e\mu/\mu e$  (top) and  $ee/\mu\mu$  (bottom) regions in the analysis. All regions are shown as counters after event selection, showing the post-fit data-to-Monte-Carlo agreement including the systematic uncertainties.

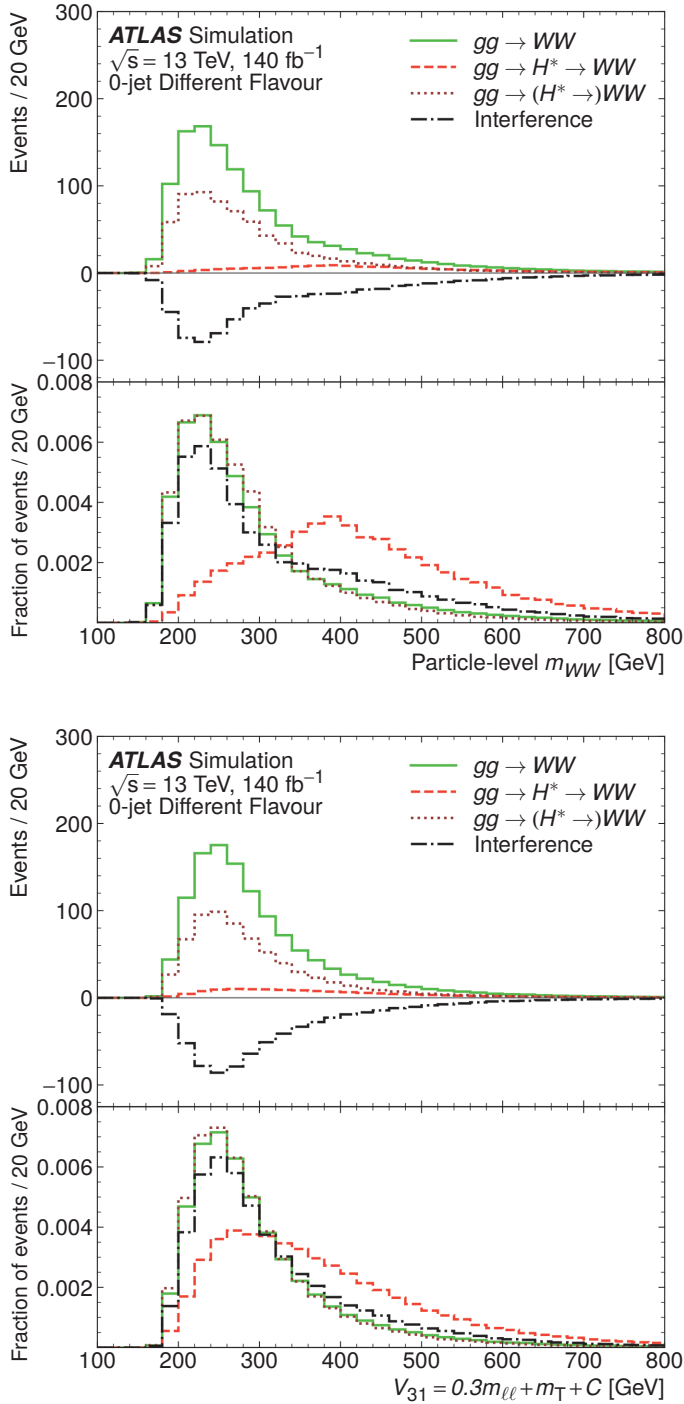
## 10.4 The $m_{WW}$ proxy

When measuring the off-shell process, it is crucial to reconstruct the invariant mass of the  $H^* \rightarrow WW$  system,  $m_{WW}$ , defined at the truth level as  $m_{WW} > 140$  GeV. Fig. 10.2 illustrates that the destructive interference, which has a shape similar to the background, is more prominent in the low- $m_{WW}$  region, while the signal is more pronounced in the high- $m_{WW}$  region. Therefore,  $m_{WW}$  serves to distinguish between low and high energetic regimes, aiding in the separation of destructive interference from the signal and enhancing the sensitivity of the measurement. However, since reconstructing  $m_{WW}$  is challenging due to the presence of neutrinos in the final state. The information provided by the charged leptons and jets is partial and limited regarding the full kinematic properties of the  $WW$  system.

To approximate  $m_{WW}$  and serve as the discriminant variable in the fit, an analytic variable,  $V_{xy}$ , is utilized. Commonly referred to as the  $m_{WW}$  proxy,  $V_{xy}$  is defined as  $V_{xy} = x \cdot m_{\ell\ell} + y \cdot m_T$ , a linear combination of  $m_{\ell\ell}$  and  $m_T$ . The observables and coefficients for  $V_{xy}$  are selected to ensure that the resulting variable closely approximates the true particle-level  $m_{WW}$  value. The choice of  $m_{\ell\ell}$  and  $m_T$  for the  $m_{WW}$  proxy is motivated by their strong correlation with the true  $m_{WW}$  at the particle-level.

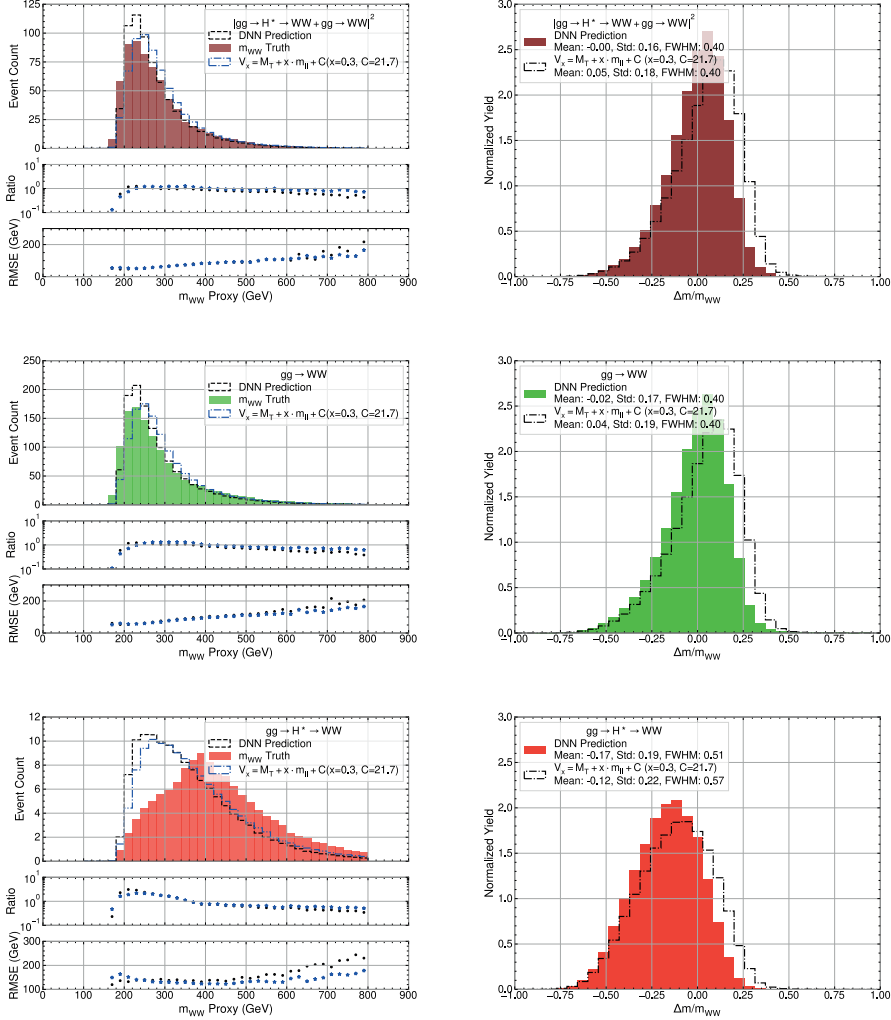
The form and coefficients of the  $V_{xy}$  variable is determined after extensive studies on its performance and comparisons with neural network regressors. Given the simplicity of  $V_{xy}$ , one might expect a more complex model, such as neural networks, to predict  $m_{WW}^{\text{truth}}$  with greater precision and accuracy, since it can utilize more kinematic variables and non-linear correlations. Therefore,  $V_{xy}$  variables are compared with an ensemble of neural network regressors. The regression target is  $m_{WW}^{\text{truth}}$ , calculated from the momentum of the final state  $\ell\nu\ell\nu$  system.

The performance of the  $m_{WW}$  proxy is evaluated across the  $N_{\text{jet}} = 0$   $e\mu/\mu e$ ,  $N_{\text{jet}} = 1$   $e\mu/\mu e$ , and  $N_{\text{jet}} \geq 2$   $e\mu/\mu e$  regions using key metrics such as mean, standard deviation, and full width at half maximum of the pull distribution, in addition to counting significance. These results are directly applicable to the  $ee/\mu\mu$  channel, as identical signal kinematics are expected in the  $ee/\mu\mu$  channel.



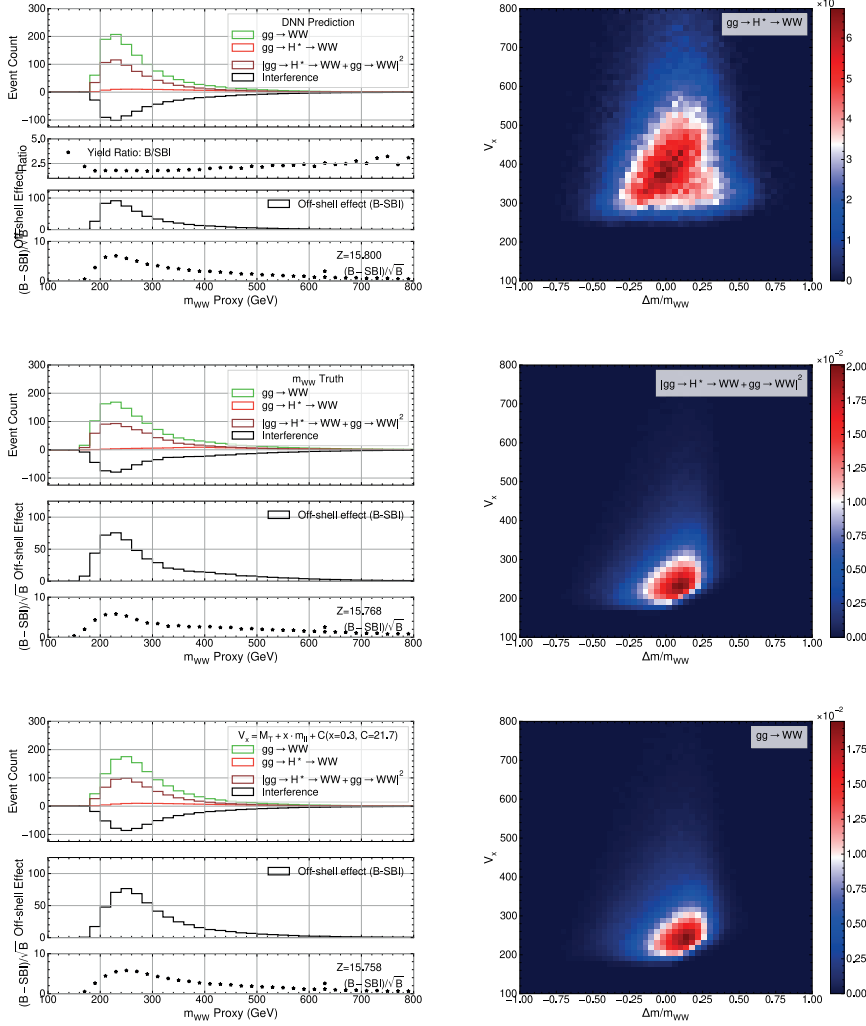
**Figure 10.2:** Distribution of the particle-level diboson invariant mass  $m_{WW}$  (top). Distribution of  $V_{0.3|1.0}$  with a constant  $C=21.7$  GeV added for an easier comparison to the particle-level diboson invariant mass. Both plots show the expected yield of  $gg \rightarrow WW$  (green, solid),  $gg \rightarrow H^* \rightarrow WW$  (red, dashed), and  $gg \rightarrow (H^*) \rightarrow WW$  (brown, dotted) events and the size of interference (black, dashed and dotted) in the  $e\mu/\mu e$  0-jet category. In the bottom pad, the histograms are normalised for comparison.

Fig. 10.3 presents the distributions of  $m_{WW}$  proxy and  $m_{WW}^{\text{truth}}$ , along with their corresponding pull plots, separately for processes including  $gg \rightarrow (H^*) \rightarrow WW$ ,  $gg \rightarrow H^* \rightarrow WW$ , and  $gg \rightarrow WW$ . The best-performing  $V_{xy}$  variable,  $V_{0.3|1.0}$ , demonstrates performance comparable to the best neural network regressors across several metrics.



**Figure 10.3:** Distributions of the DNN regressor,  $V_{0.3|1.0}$ , and  $m_{WW}^{\text{truth}}$  for  $N_{\text{jet}} = 0$   $\mu\mu/\mu e$  on the left. On the right, the corresponding difference between the truth and the predicted value by the DNN regressor and  $V_{0.3|1.0}$  are shown. The plots show the distributions and pull plots for the  $gg \rightarrow (H^*) \rightarrow WW$ ,  $gg \rightarrow WW$  and  $gg \rightarrow H^* \rightarrow WW$  processes from top to bottom.

Additionally, the predicted distributions and counting significance of the  $m_{WW}$  proxy and  $m_{WW}^{\text{truth}}$  are shown in Fig. 10.4. The difference between the  $\mu_{\text{off-shell}} = 1$  and  $\mu_{\text{off-shell}} = 0$  hypotheses is referred to as the "off-shell effect" which forms the basis for calculating the counting significance.



**Figure 10.4:** Distributions of the DNN regressor,  $V_{0.3|1.0}$ , and  $m_{WW}^{\text{truth}}$  for  $N_{\text{jet}} = 0$   $e\mu/\mu e$  on the left row from top to bottom. The panels show the ratio of interfering background to off-shell effect, the residual of the off-shell effect and the counting significance of the off-shell effect with respect to the interfering background. On the right, from top to bottom, the difference between  $m_{WW}^{\text{truth}}$  and the predicted value by the chosen  $m_{WW}$  proxy,  $V_{0.3|1.0}$  are shown for the S, SBI and B processes.

The counting significances of the DNN,  $V_{xy}$  variables, and  $m_{WW}^{\text{truth}}$  are found to be very close. These studies conclude that the performance of  $V_{xy}$  variables is comparable to that of the best-performing regression neural networks. Despite its simplicity,  $V_{0.3|1.0}$  achieves a counting significance close to that of  $m_{WW}^{\text{truth}}$ , making it an excellent  $m_{WW}$  proxy.



## 10.5 Machine learning strategy

Dense neural networks are trained using Tensorflow [181] for each signal region to distinguish the signal ( $gg \rightarrow H^* \rightarrow WW \rightarrow \ell\nu\ell\nu$ ,  $qq \rightarrow H^* \rightarrow WW + jj \rightarrow \ell\nu\ell\nu + jj$ ) and its interfering background from the non-interfering background. This enhances the off-shell effect and simplifies the optimization to a binary classification in each region.

The DNN target in each region consists of the combined S and SBI samples of the main signal contributions, such as ggF S and SBI for the  $N_{\text{jet}} = 0$  region. For the  $N_{\text{jet}} = 1$  region, due to a higher number of VBF events, VBF S and SBI is also included along with ggF S and SBI. In the  $N_{\text{jet}} \geq 2$  DNN region, the target is VBF S and SBI. The usage of both SBI and S as target serves multiple purposes: it prevents discrimination against interfering backgrounds (and thus interference), stabilizes the training with the larger number of events in the full SBI sample, particularly for VBF, and does not significantly reduce the discriminating power of S to background compared to using S alone. At last, all the non-interfering backgrounds are grouped under the non-target label.

A standardized protocol is developed during the analysis to formalize the approach for each signal region. Training and binning optimization is performed in relatively inclusive signal regions in order to maximize the impact of the DNN's. First, training samples are created by producing ROOT files which contain MC events from each signal region. Instead of data-driven fakes, the W+jets Monte Carlo sample are used to mimic non-prompt lepton contributions. To avoid biases, a 5-fold cross-validation is applied to the dataset, which is split into training (3 folds), validation (1 fold), and test (1 fold) samples. A preprocessing routine balances the training set and applies symmetries (explained below) to reduce input variables. The set of input variables is minimized further by examining each variable's impact on the model prediction. The best model is achieved through hyperparameter scanning and selecting the model with the best performance metrics. The process is summarized in the following steps:

- **reweighting:** Applying weights to the targets (SBI and S) such that their sum equals that of the background events.
- **Rotation:** Rotation of the leading-lepton  $\phi$  coordinate to 0 and adjusting  $\phi$  values of other objects accordingly, which reduces the number of training input variables by 1, since the  $\phi$  of the leading-lepton is always 0.
- **Mirroring:** Projecting all events'  $\eta$  variables along the xy plane, based on a chosen object's  $\eta$ , exploiting the event symmetry. In DNN training, adjusting  $\eta$  values after rotation ensures consistency.
- **Hyperparameter scan:** Searching for optimal hyperparameters across a comprehensive set of observables expected to enhance model performance. The best model, shown in Tab. 10.2, based on Poisson significance and ROC AUC metrics, is selected from approximately 300 scans.

- **Permutation:** Iteratively shuffling input observables to gauge their impact on the model performance. Significant changes (arbitrarily defined as 1% deviation) flag relevant observables for subsequent trainings.
- **Reduction due to mis-modelling:** Plotting reduced observables in control regions to identify and exclude poorly modelled variables (e.g., jet count affected by parton showering).
- **Retrain:** Retraining the best performing model with the reduced set of input variables.

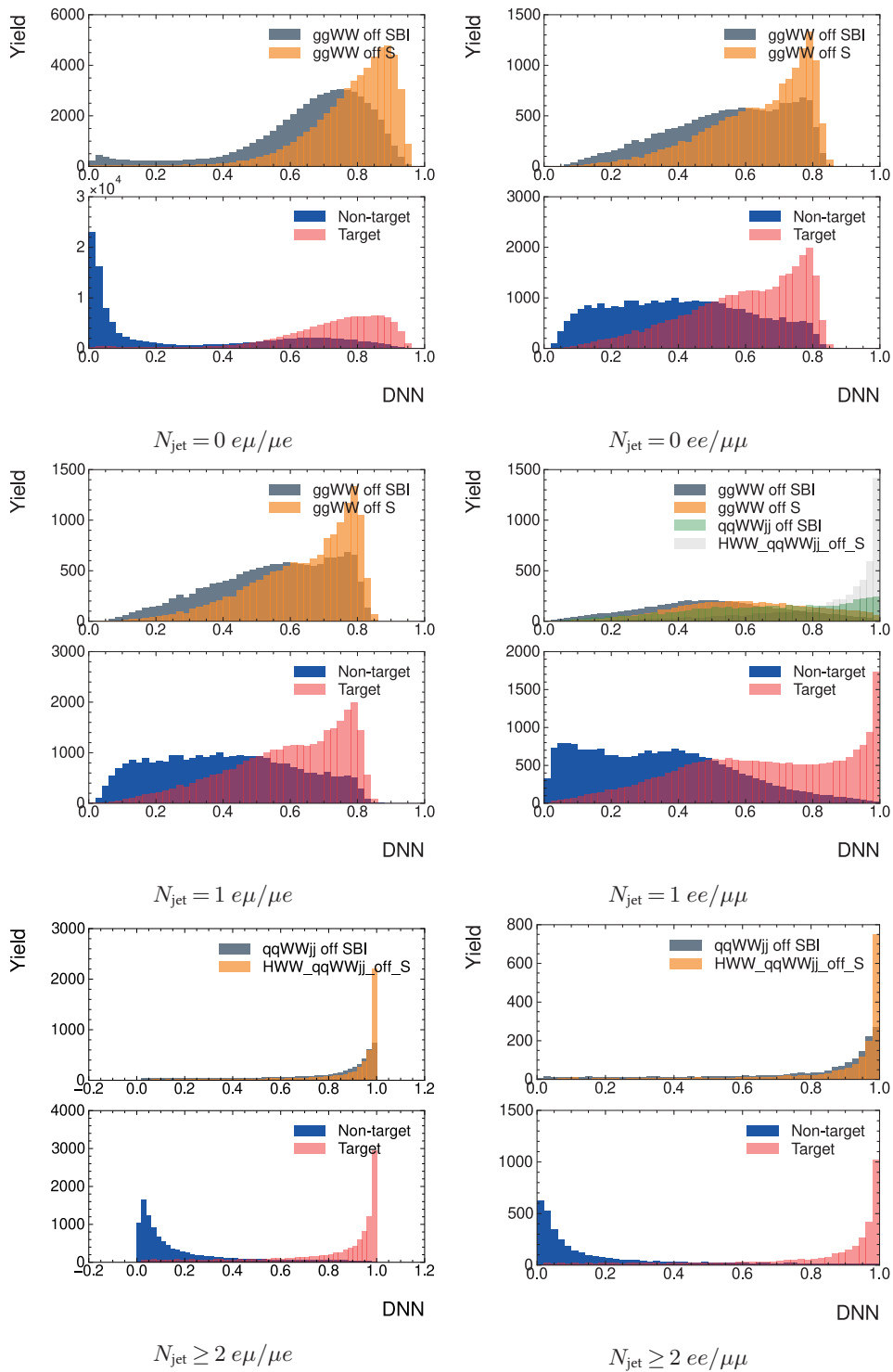
The final set of the input observables for each region is summarized in [Tab. 10.3](#), where the check marks show whether the variable is used in a region or not. The resulting discriminant variables for each of the jet multiplicity regions can be seen in [Fig. 10.5](#). The strength of the input variables indicated by the permutation test can be visualized by illustrating the optimization target versus the background in 1D observable plots in [Fig. 10.6](#).

**Table 10.2:** Hyperparameters for the neural network trained for each signal region. The activation functions `relu` and `elu` stand for the rectified linear unit and exponential linear unit activation functions, respectively. The optimizer used is the stochastic gradient descent (SGD) algorithm. The learning rate decay ratio is the factor by which the learning rate is multiplied after each epoch. The Nesterov momentum is used in the  $ee/\mu\mu$  1-jet region, which improves the convergence speed of SGD. The dropout is the fraction of the input units to drop. The L2 regularization weight is the factor by which the L2 regularization term is multiplied.

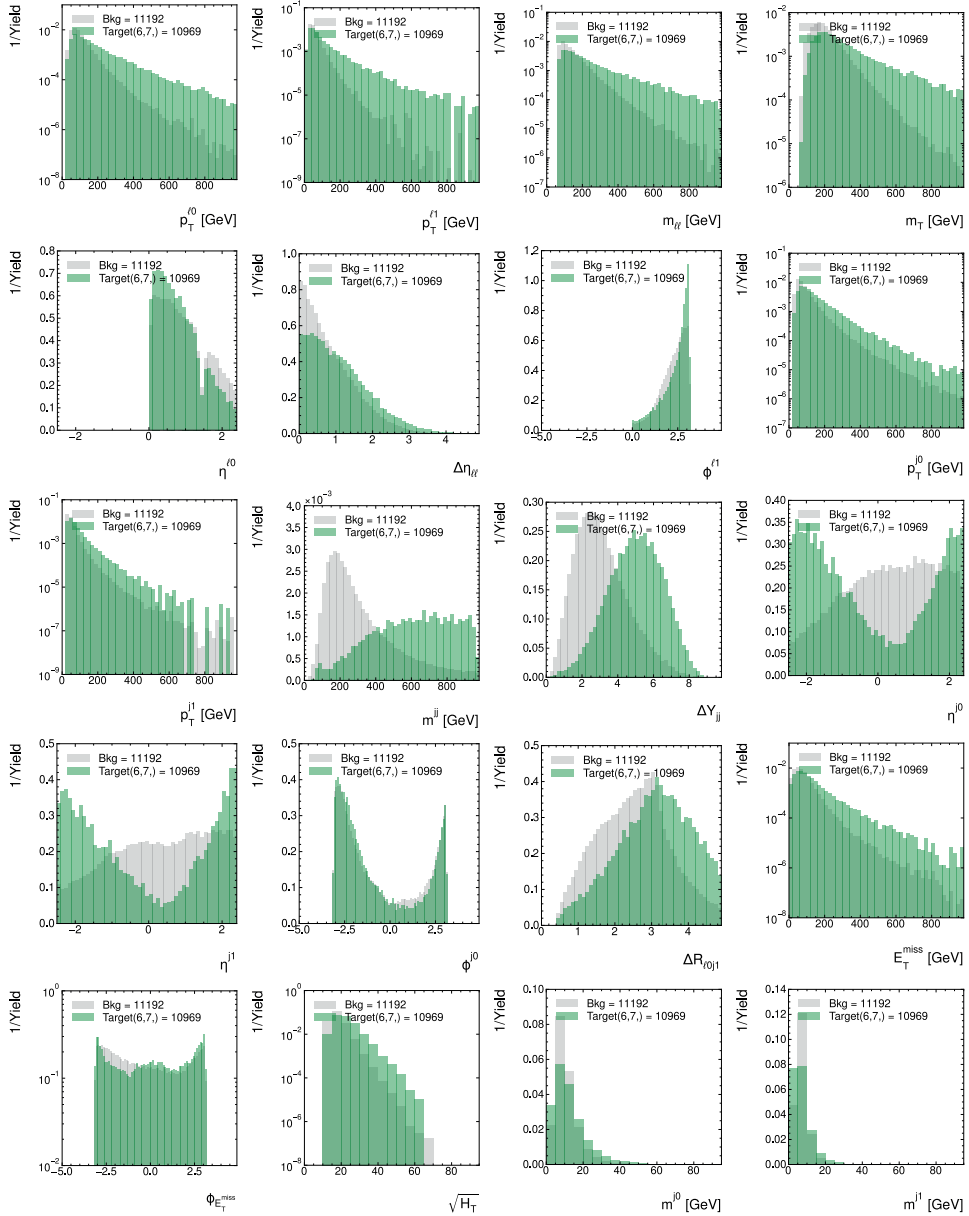
| Hyperparameter            | $e\mu/\mu e$ leptons |       |               | $ee/\mu\mu$ leptons |       |               |
|---------------------------|----------------------|-------|---------------|---------------------|-------|---------------|
|                           | 0-jet                | 1-jet | $\geq 2$ -jet | 0-jet               | 1-jet | $\geq 2$ -jet |
| Nodes in layer 1          | 256                  | 128   | 256           | 256                 | 256   | 128           |
| Nodes in layer 2          | 128                  | 32    | 128           | 128                 | 64    | 128           |
| Nodes in layer 3          | 64                   | 32    | 64            | 64                  | 256   | 32            |
| Nodes in layer 4          | 16                   |       | 16            | 16                  |       | 16            |
| Nodes in layer 5          | 16                   |       | 16            |                     |       |               |
| Activation function       | relu                 | relu  | elu           | relu                | elu   | relu          |
| optimizer                 | SGD                  | SGD   | SGD           | SGD                 | SGD   | SGD           |
| Epochs                    | 300                  | 300   | 30            | 300                 | 100   | 300           |
| Batch-size                | 256                  | 512   | 128           | 256                 | 1024  | 256           |
| Learning rate             | 0.3                  | 0.1   | 0.1           | 0.1                 | 0.1   | 0.03          |
| Learning rate decay ratio | 0.01                 | 0.8   | 0.01          | 0.3                 | 1     | 1             |
| Use Nesterov momentum?    |                      |       |               |                     | ✓     |               |
| Momentum                  |                      |       |               |                     | 0.95  |               |
| Dropout                   | 0.3                  | 0.1   | 0.6           | 0.6                 | 0     | 0             |
| L2 regularization weight  | 0                    | 0     | 0.0003        | 0.0003              | 0     | 0.00003       |

**Table 10.3:** Variables used as neural network inputs in each of the analysis SRs.

| Variable                   | different-flavour lepton category |       |               | same-flavour lepton category |       |               |
|----------------------------|-----------------------------------|-------|---------------|------------------------------|-------|---------------|
|                            | 0-jet                             | 1-jet | $\geq 2$ -jet | 0-jet                        | 1-jet | $\geq 2$ -jet |
| $p_T^{\ell 0}$             | ✓                                 | ✓     | ✓             | ✓                            | ✓     |               |
| $p_T^{\ell 1}$             | ✓                                 | ✓     | ✓             | ✓                            | ✓     | ✓             |
| $\eta^{\ell 0}$            | ✓                                 | ✓     | ✓             | ✓                            | ✓     | ✓             |
| $\eta^{\ell 1}$            | ✓                                 | ✓     | ✓             | ✓                            | ✓     | ✓             |
| $\phi^{\ell 0}$            | ✓                                 | ✓     | ✓             | ✓                            | ✓     | ✓             |
| $\phi^{\ell 1}$            | ✓                                 |       | ✓             | ✓                            | ✓     | ✓             |
| $p_T^{\ell\ell}$           |                                   | ✓     |               |                              |       |               |
| $\Delta\eta^{\ell\ell}$    |                                   | ✓     | ✓             |                              |       | ✓             |
| $\Delta y^{\ell\ell}$      |                                   | ✓     | ✓             |                              | ✓     | ✓             |
| $\Delta\phi^{\ell\ell}$    |                                   |       | ✓             |                              | ✓     | ✓             |
| $\Delta R^{\ell\ell}$      |                                   |       |               |                              | ✓     |               |
| $m_T$                      | ✓                                 | ✓     | ✓             | ✓                            | ✓     | ✓             |
| $m_{\ell\ell}$             | ✓                                 | ✓     | ✓             | ✓                            | ✓     | ✓             |
| $\max(m_T^W)$              | ✓                                 |       |               | ✓                            | ✓     |               |
| $p_T^{j0}$                 |                                   | ✓     | ✓             |                              | ✓     | ✓             |
| $p_T^{j1}$                 |                                   |       | ✓             |                              |       | ✓             |
| $\eta^{j0}$                |                                   | ✓     | ✓             |                              | ✓     | ✓             |
| $\eta^{j1}$                |                                   |       | ✓             |                              |       | ✓             |
| $\phi^{j0}$                |                                   |       | ✓             |                              | ✓     | ✓             |
| $m^{j0}$                   |                                   | ✓     | ✓             |                              | ✓     | ✓             |
| $m^{j1}$                   |                                   |       | ✓             |                              |       | ✓             |
| $m^{jj}$                   |                                   |       | ✓             |                              |       | ✓             |
| $\Delta y^{jj}$            |                                   |       | ✓             |                              |       | ✓             |
| $\sqrt{H_T}$               |                                   | ✓     | ✓             |                              |       | ✓             |
| $E_T^{\text{miss}}$        | ✓                                 | ✓     | ✓             | ✓                            |       | ✓             |
| $\phi_{p_T^{\text{miss}}}$ | ✓                                 |       | ✓             | ✓                            |       | ✓             |
| $S_{\text{miss}}$          | ✓                                 |       | ✓             | ✓                            | ✓     | ✓             |
| $\Delta R^{\ell 0 j 0}$    |                                   | ✓     |               |                              | ✓     |               |
| $\Delta R^{\ell 0 j 1}$    |                                   |       | ✓             |                              |       | ✓             |
| $\Delta R^{\ell 1 j 0}$    |                                   | ✓     |               |                              | ✓     |               |



**Figure 10.5:** The DNN discriminant variable for each signal region. In the top pad the distribution of the selected DNN targets is shown. The bottom pad shows the distribution of the combined target (interfering processes) against all non-target (every other non-interfering background) processes.

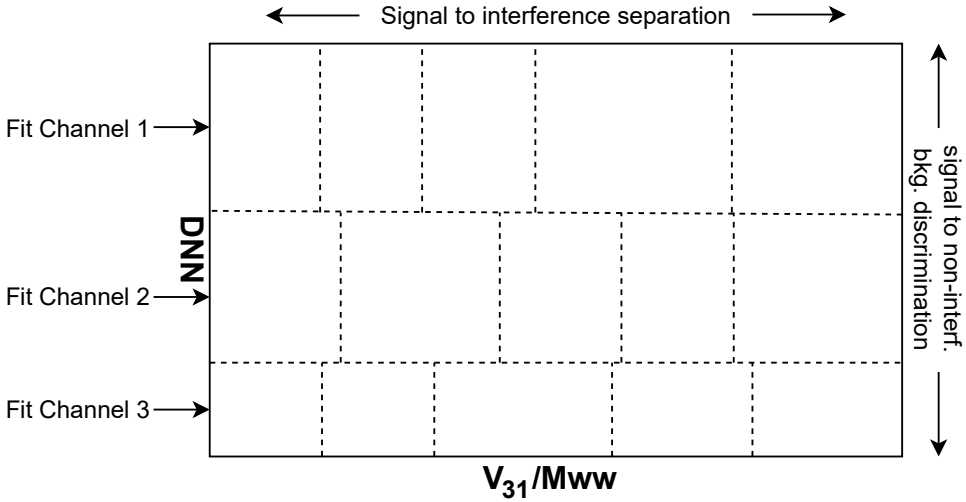


**Figure 10.6:** A selection of input DNN distributions with a comparison between target ( $qq \rightarrow (H^*) \rightarrow WW + jj \rightarrow \ell\nu\ell\nu + jj$  S and SBI (labelled 6 and 7 in the optimization framework)) in green and background (non-target) in grey (every other background) in the  $N_{\text{jet}} \geq 2$   $e\mu/\mu e$  region. The total number of simulated events for each process is also shown in the legend.

## 10.6 Unrolling the DNN

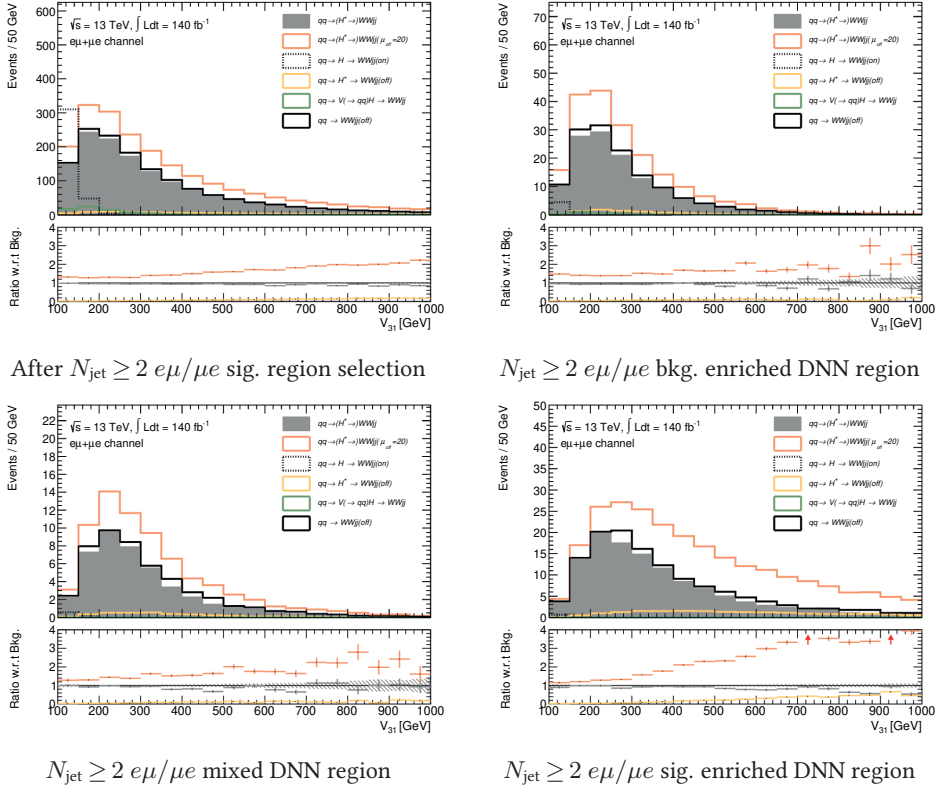
In addition to separating signal from background, DNN discriminants are commonly used as final fitting variables due to their sensitivity to the target. In this analysis, the signal-background separation is handled by the discriminant, fitting is conducted on  $V_{0.3|1.0}$  as detailed in § 10.4. However,  $V_{0.3|1.0}$  provides minimal separation between off-shell effects and the non-interfering background.

In order to optimize the signal purity without event loss, an unrolling approach divides the DNN distribution into three regions, which enables a 2D fit per region as depicted in Fig. 10.7. Via the use of this 2D fit, the analysis combines the background-separation capabilities of the DNNs with the signal-interference separation capabilities of the  $V_{0.3|1.0}$  variable.



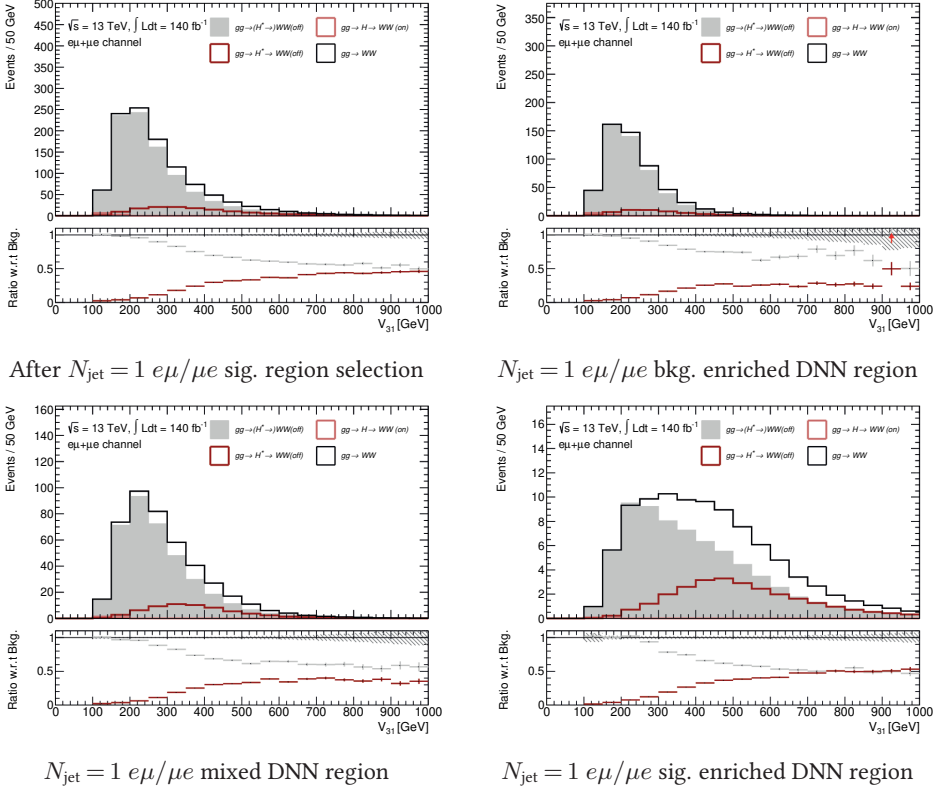
**Figure 10.7:** An illustration of the optimization strategy in the off-shell analysis.

The final  $m_{WW}$  proxy distributions with event selections applied, both pre- and post-DNN unrolling, for the signal templates in this analysis, are shown in Figs. 10.8 and 10.9 for the  $N_{\text{jet}} = 1$   $e\mu/\mu e$  and  $N_{\text{jet}} \geq 2$   $e\mu/\mu e$  regions. These figures illustrate the increase in the off-shell effect in the bulk of the  $V_{0.3|1.0}$  distribution, as well as in the tail in each of the signal regions, when the DNN unrolling is applied. It can be seen that the ratio of the signal and SBI samples with respect to the interfering background increases and decreases, respectively. The  $ee/\mu\mu$  analysis regions are not shown but follow a similar trend.



**Figure 10.8:** Unrolled  $m_{WW}$  proxy ( $V_{0.3|1.0}$ ) signal distributions and their ratio w.r.t. the interfering background with statistical uncertainties. The grey and black histograms are the SBI contribution (includes the interference) and the background-only distribution (excludes the interference), respectively. The striped and yellow histograms show the on-shell and off-shell VBF signal contributions, respectively. The orange histogram shows the VBF SBI sample with  $\mu_{\text{off-shell}} = 20$  and the green histogram shows the on-shell  $VH$  contribution which ends up in the signal region.

A simultaneous binned profile likelihood fit to  $V_{0.3|1.0}$  is conducted in each DNN region, demonstrating superior performance over simple  $V_{0.3|1.0}$  fitting in both the  $N_{\text{jet}} = 0$   $e\mu/\mu e$  and  $N_{\text{jet}} \geq 2$   $e\mu/\mu e$  channels. This method is particularly effective at the  $\mu_{\text{off-shell}}^{\text{ggF0j}} = 0$  exclusion limit. However, the strategy of including the mixed and background-enriched DNN regions in the fit is less impactful in the  $N_{\text{jet}} \geq 2$  channel, where the DNN performs more effectively compared to the  $N_{\text{jet}} = 0$  and  $N_{\text{jet}} = 1$  channels.



**Figure 10.9:** Unrolled  $m_{WW}$  proxy ( $V_{0.3|1.0}$ ) signal distributions and their ratio w.r.t. the interfering background with statistical uncertainties. Grey and black histograms are the SBI contribution (includes the interference) and the background-only distribution (excludes the interference), respectively. Light and dark red histograms show the on-shell and off-shell ggF signal, respectively.

### 10.6.1 Binning optimization

Several binning criteria are applied to determine the optimal  $V_{0.3|1.0}$  binning, listed as follows:

- The number of signal events in each bin must be equal to or greater than 1, aligned with the DNN optimization targets discussed previously. Tests with different signal definitions showed minimal impact on bin boundaries.
- A minimum of 10 non-interfering background events per bin is required to estimate theory uncertainties reliably and avoid empty bins in the data distribution.
- Statistical uncertainty on signal and background events has to be less than 20%.
- A hard lower limit of 10 GeV is set for bin sizes in order to avoid a scenario where the number of bins causes numerical issues in the fit.
- The D-transformation, commonly used in top and Higgs analyses, is applied to reduce the number of bins for stability and convergence of the fit. This criterion



adjusts the binning dynamically using weights  $w_s$  and  $w_b$ , ensuring fewer bins in lower DNN regions and more in higher DNN regions. It is based on the  $Z$ -value, defined as:

$$Z = w_s \frac{\sum_{i=1}^k n_s^i}{N_s} + w_b \frac{\sum_{i=1}^k n_b^i}{N_b},$$

where  $N_s$  and  $N_b$  are the total numbers of signal and background events, respectively. The numerator represents the sum of signal and background yields from bin  $i$  to bin  $k$ . The algorithm begins with fine binning, starting from bin  $i = n_{\text{bins}}$  and progressively merges nearby bins in reverse order. Each time bins are merged, a new value of  $Z$  is computed. The sum in the numerator continues as long as  $Z \leq 1$ . Once  $Z > 1$ , a new coarser bin is created with its new edge defined by the current bin  $j$ . Then,  $Z$  is reset to 0,  $i$  is set to  $j$ , and the algorithm proceeds again until  $Z > 1$ . Here,  $w_s$  and  $w_b$  are freely tunable weights to adjust the binning granularity. Specifically,  $w_s$  and  $w_b$  are chosen as:  $w_s = 0$  and  $w_b = 5$  for the background-enriched DNN region,  $w_s = 10$  and  $w_b = 10$  for the mixed region, and  $w_s = 10$  and  $w_b = 0$  for the signal-enriched region. These choices maximize sensitivity, without compromising the fit stability due to a large number of bins.

### 10.6.2 DNN boundary optimization

The previous subsection demonstrated the optimization of the  $V_{0.3|1.0}$  variable. In its application in this analysis, adjusting the boundaries of the DNN distribution can influence the signal purity and final results. Consequently, a scan of the DNN boundaries is conducted, where both boundaries are varied. Subsequent unrolling and likelihood fits are performed to determine new significances, as well as  $1\sigma$  and  $2\sigma$  values for the POI  $\mu_{\text{off-shell}}$ . By applying simultaneous optimization of  $V_{0.3|1.0}$  and DNN, the optimal DNN boundaries shown in [Tab. 10.4](#) are determined.

**Table 10.4:** Optimal DNN boundaries for each measured region, as determined by the optimization procedure.

| Channel                                     | Background-enriched upper bound | Mixed upper bound |
|---------------------------------------------|---------------------------------|-------------------|
| $N_{\text{jet}} = 0 \text{ } e\mu/\mu e$    | 0.61                            | 0.82              |
| $N_{\text{jet}} = 1 \text{ } e\mu/\mu e$    | 0.47                            | 0.73              |
| $N_{\text{jet}} \geq 2 \text{ } e\mu/\mu e$ | 0.82                            | 0.91              |
| $N_{\text{jet}} = 0 \text{ } ee/\mu\mu$     | 0.64                            | 0.78              |
| $N_{\text{jet}} = 1 \text{ } ee/\mu\mu$     | 0.62                            | 0.79              |
| $N_{\text{jet}} \geq 2 \text{ } ee/\mu\mu$  | 0.67                            | 0.95              |

The fit includes simplified statistical Monte Carlo uncertainties and systematic uncertainties that affect the overall normalization of a process. Overall, the fits reveal no dominant bin leading the sensitivity, but rather an area of sensitivity where differences among neighbouring bins are less than 5%, suggesting a certain flexibility in the choice.

## 10.7 Control and validation regions

When a signal region is defined as a phase-space region enriched in signal events, a proper estimation of the background events in that signal region is needed in order to measure the signal yield with enough accuracy. In order to ensure the Monte Carlo simulation properly predicts a given background process, a modelling check is performed in a region that is as close as possible to the signal region. In this region, the background of interest should be dominant and is compared to the data.

Given the number of observed events in the control region,  $N_{\text{CR}}^{\text{data}}$ , the normalization factor is defined as the ratio of the number of observed events in the control region to the number of expected background events in the control region,  $N_{\text{CR}}^{\text{MC,bkg}}$ . The normalization factor,

$$\text{NF} = \frac{N_{\text{CR}}^{\text{data}}}{N_{\text{CR}}^{\text{MC,bkg}}},$$

can then be applied to the expected number of background events in the control region to improve agreement between data and Monte Carlo. Additionally, the NF calculated from the CR is applied to the background events in the signal region,  $N_{\text{SR}}^{\text{MC,bkg}}$ , to obtain a better estimate of the expected number of background events in the signal region,  $N_{\text{SR}}^{\text{MC,bkg}}$ . By calculating the expected background events in the signal region as  $N_{\text{SR}}^{\text{MC,corr,bkg}} = \text{NF} \cdot N_{\text{SR}}^{\text{MC,bkg}}$ , a more accurate integrated background yield,  $N_{\text{SR}}^{\text{MC,corr,bkg}}$ , can be estimated without affecting any kinematic observables in the signal region.

This normalization, however, is extrapolated from the control region. This makes it important that the extrapolation uncertainties between the control region and signal region must be as small as possible in order to make the extrapolation reliable.

The off-shell analysis uses multiple control regions that control the normalization factors of many backgrounds. Therefore, in the final profile likelihood fit, the normalization factors are included as nuisance parameters. This fully takes into account correlations, and accurately describes the yield in the signal and control regions depending on the fit convergence quality. On top of the post-fit normalization factors (NFs), a good estimate of the background modelling is also needed before the application of the statistical treatment and uncertainties. Therefore, pre-fit NFs are estimated using the matrix inversion method.

This method allows for the estimation of the normalization factors in the control regions via a matrix inversion, taking into account the correlations between the backgrounds in the various control regions. The problem of calculating two or more NFs from, for example, the top,  $Z \rightarrow \ell\ell$  or more control regions can be cast in matrix notation

$$\begin{pmatrix} N_{\text{data}}^{\text{CR,top}} \\ \vdots \\ N_{\text{data}}^{\text{CR},Z \rightarrow \ell\ell} \end{pmatrix} = \begin{pmatrix} N_{\text{MC,top}}^{\text{CR,top}} & \cdots & N_{\text{MC},Z \rightarrow \ell\ell}^{\text{CR,top}} \\ \vdots & \ddots & \vdots \\ N_{\text{MC,top}}^{\text{CR},Z \rightarrow \ell\ell} & \cdots & N_{\text{MC},Z \rightarrow \ell\ell}^{\text{CR},Z \rightarrow \ell\ell} \end{pmatrix} \cdot \begin{pmatrix} \text{NF}_{\text{top}} \\ \vdots \\ \text{NF}_{Z \rightarrow \ell\ell} \end{pmatrix},$$

and then solved for the NFs by inverting the matrix

$$\begin{pmatrix} \text{NF}_{\text{top}} \\ \vdots \\ \text{NF}_{Z \rightarrow \ell\ell} \end{pmatrix} = \begin{pmatrix} N_{\text{MC,top}}^{\text{CR,top}} & \cdots & N_{\text{MC,Z} \rightarrow \ell\ell}^{\text{CR,top}} \\ \vdots & \ddots & \vdots \\ N_{\text{MC,top}}^{\text{CR,Z} \rightarrow \ell\ell} & \cdots & N_{\text{MC,Z} \rightarrow \ell\ell}^{\text{CR,Z} \rightarrow \ell\ell} \end{pmatrix}^{-1} \cdot \begin{pmatrix} N_{\text{data}}^{\text{CR,top}} \\ \vdots \\ N_{\text{data}}^{\text{CR,Z} \rightarrow \ell\ell} \end{pmatrix}.$$

The result of this matrix can then be used to calculate the expected number of background events in the signal region,  $N_{\text{SR}}^{\text{MC,corr,bkg}}$ , by the matrix equation

$$N_{\text{SR}}^{\text{MC,corr,bkg}} = \begin{pmatrix} N_{\text{SR}}^{\text{top,MC}} & \cdots & N_{\text{SR}}^{\text{Z} \rightarrow \ell\ell, \text{MC}} \end{pmatrix} \cdot \begin{pmatrix} \text{NF}_{\text{top}} \\ \vdots \\ \text{NF}_{Z \rightarrow \ell\ell} \end{pmatrix},$$

in order to get a reliable indication of the overall background normalization in each control and signal region, before any fits with uncertainties have taken place. In the following sections, the plots shown for the control and validation regions are normalized using the matrix inversion method.

### 10.7.1 Event selection of the control and validation regions

The control and validation region selection criteria are defined in [Tab. 10.5](#), in addition to the  $e\mu/\mu e$  and  $ee/\mu\mu$  pre-selections summarized in [Tab. 10.1](#). All control regions are designed such that they are enriched with the target background and orthogonal to the off-shell signal regions, the on-shell analysis and the  $H^* \rightarrow ZZ$  off-shell measurement.

Although the control regions are used inclusively in the fit, the extracted normalization factors are used to normalize the three DNN regions simultaneously, while unrolled in the  $V_{0.3|1.0}$  distribution. Therefore, it is important to also validate the modelling in this configuration and make sure the data agrees with the Monte Carlo within uncertainty.

### 10.7.2 The $qq \rightarrow WW$ control and validation regions

The  $qq \rightarrow WW$  process represents the dominant background in the  $N_{\text{jet}}=0$   $e\mu/\mu e$  region with its contribution reaching about 65% of the total background. A validation region is defined in the  $N_{\text{jet}}=0$   $e\mu/\mu e$  channel to cross-check the reliability of the modelling in this channel. The  $qq \rightarrow WW$  validation region is defined by the same selection criteria as the signal region, except for the  $\Delta\eta(\ell, \ell)$  requirement, which is inverted to  $\Delta\eta(\ell, \ell) > 1.8$ .

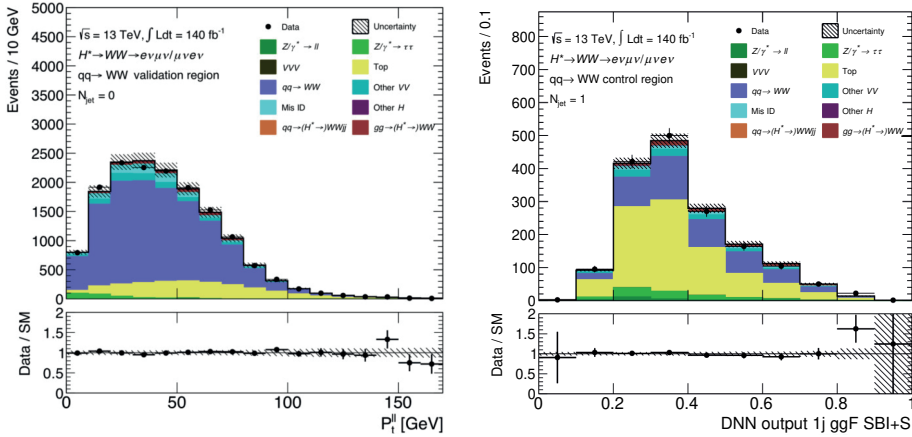
The  $N_{\text{jet}}=1$   $e\mu/\mu e$  channel also contains sizeable contributions from the  $qq \rightarrow WW$  background, with the  $qq \rightarrow WW$  process contributing about 40% of the total background. An  $N_{\text{jet}}=1$   $e\mu/\mu e$  control region is designed in order to validate and control the normalization of this background in the fit.

**Table 10.5:** Event selection criteria for the validation and control regions, defined in addition to the pre-selections summarized in Tab. 10.1. The  $Z \rightarrow \ell\ell$  regions are defined using the  $ee/\mu\mu$  lepton pre-selection and categorisation requirements, the remaining regions are defined using the  $e\mu/\mu e$  requirements.

|                                            | 0-jet                                                                                                                                                                            | 1-jet | $\geq 2$ -jet                                 |
|--------------------------------------------|----------------------------------------------------------------------------------------------------------------------------------------------------------------------------------|-------|-----------------------------------------------|
| $e\mu/\mu e$ lepton category pre-selection |                                                                                                                                                                                  |       |                                               |
| $qq \rightarrow WW$                        | $n_{b\text{-jets}} = 0$<br>$m_{\ell\ell} > 100 \text{ GeV}$<br>$\Delta\eta(\ell, \ell) > 1.8$<br>$ m_{\tau\tau} - 91\text{GeV}  > 25 \text{ GeV}$                                |       | -                                             |
| Top quark                                  | $n_{b\text{-jets}} = 1$<br>$m_{\ell\ell} > 100 \text{ GeV}$<br>$\Delta\eta(\ell, \ell) < 1.8$<br>$\Delta\phi(\ell, \ell) > 1.8$                                                  |       | $m_{\ell\ell} > 70 \text{ GeV}$<br>CJV<br>OLV |
| $Z \rightarrow \tau\tau$                   | $n_{b\text{-jets}} = 0$<br>$m_{\ell\ell} < 80 \text{ GeV}$<br>$\Delta\phi(\ell, \ell) > 2.4$<br>$m_{\tau\tau} > 66 \text{ GeV}$                                                  |       | -                                             |
| $ee/\mu\mu$ lepton category pre-selection  |                                                                                                                                                                                  |       |                                               |
| $Z \rightarrow \ell\ell$                   | $n_{b\text{-jets}} = 0$<br>$\Delta R(\ell, \ell) > 1.8$<br>$ m_{\ell\ell} - 91\text{GeV}  < 15 \text{ GeV}$<br>$m_{\ell\ell} > 55 \text{ GeV}$<br>$\Delta\phi(\ell, \ell) > 1.8$ |       | $m_{\ell\ell} > 70 \text{ GeV}$<br>CJV<br>OLV |

By the requirement of  $\Delta\phi(\ell, \ell) < 1.8$ , the  $qq \rightarrow WW$  control region is orthogonal to the signal region. With an additional selection of  $|m_{\tau\tau} - 91\text{GeV}| > 25 \text{ GeV}$ , the  $Z \rightarrow \tau\tau$  background is suppressed and made orthogonal to the  $Z \rightarrow \tau\tau$  control region. The  $N_{\text{jet}} \geq 2$   $e\mu/\mu e$  channel contains a smaller contribution from the  $qq \rightarrow WW$  background and can be controlled via the background-enriched DNN regions.

In all regions, the  $qq \rightarrow WW$  background normalization is freely floating in the fit, where it is mainly constrained via dedicated control regions or background-enriched DNN regions. The inclusive  $p_T^{\ell\ell}$  in the  $N_{\text{jet}} = 0$   $qq \rightarrow WW$  validation and the output of the  $N_{\text{jet}} = 0$   $e\mu/\mu e$  DNN in the  $N_{\text{jet}} = 1$   $qq \rightarrow WW$  control region are illustrated in Fig. 10.10 and show good agreement with data.



**Figure 10.10:** The  $p_T^{\ell\ell}$  in the  $N_{\text{jet}} = 0$   $qq \rightarrow WW$  validation region and the output of the  $N_{\text{jet}} = 0$   $e\mu/\mu e$  DNN in the  $N_{\text{jet}} = 1$   $qq \rightarrow WW$  the control region. The uncertainty corresponds to the total theoretical, experimental and statistical uncertainty.

### 10.7.3 The top control regions

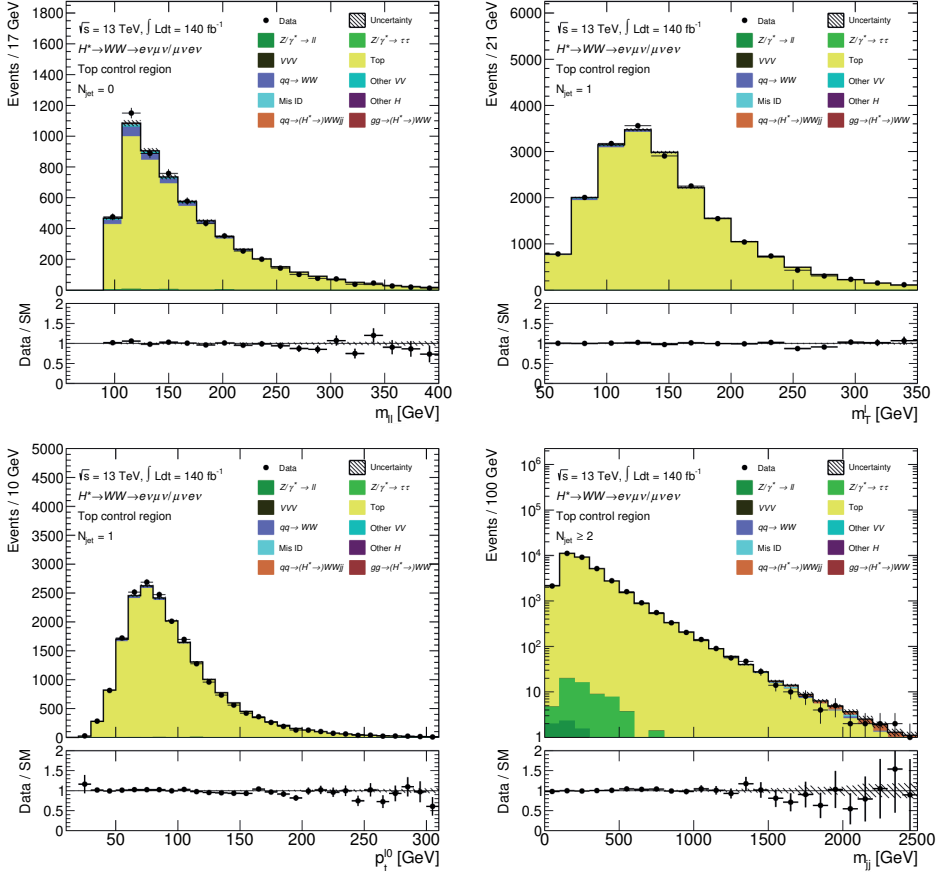
Except for the  $N_{\text{jet}} = 0$   $e\mu/\mu e$  channel, the top background is the next largest background in all of the different-flavour channels. Inclusive distributions of the  $V_{0.3|1.0}$  variable in the top control regions are shown in Fig. 10.11. In total, three top control regions are setup in each of the  $e\mu/\mu e$   $N_{\text{jet}}$  channels, where the region follows the exact same selection as the corresponding  $N_{\text{jet}}$  signal region, except for an inverted  $b$ -jet veto and a  $m_{\ell\ell} > 100$  GeV requirement across the  $N_{\text{jet}} = 0$  and  $N_{\text{jet}} = 1$  region and a  $m_{\ell\ell} > 70$  GeV requirement in the  $N_{\text{jet}} \geq 2$  region.

The top control regions are designed to control the normalization of the top background in the fit in both  $e\mu/\mu e$  and  $ee/\mu\mu$  channels. In the  $N_{\text{jet}} = 0$  channel, the tagged  $b$ -jet is required to have  $20 < p_T < 30$  GeV in order to be as close as possible to the  $b$ -jets entering the  $N_{\text{jet}} = 0$  signal region. In the  $N_{\text{jet}} = 1$  and  $N_{\text{jet}} \geq 2$  channel, the  $b$ -jet is required to have  $p_T > 30$  GeV. The top background normalization is left floating freely in the fit and is constrained via the control regions.

### 10.7.4 The $Z \rightarrow \tau\tau$ control and validation regions

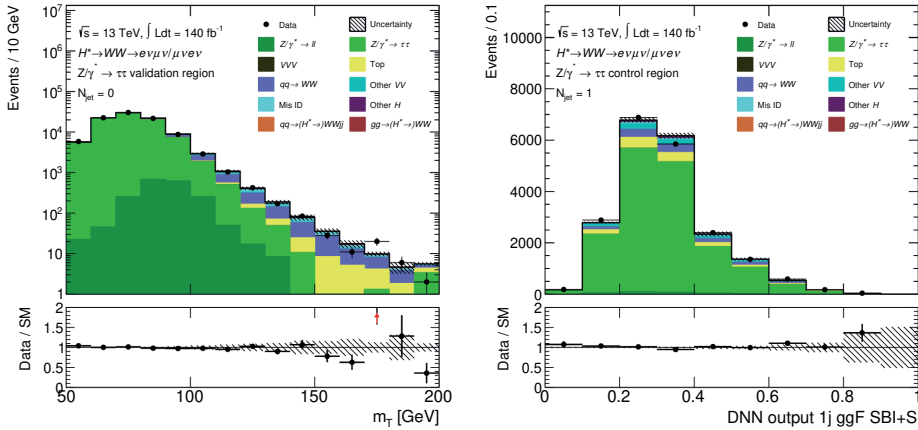
A  $Z \rightarrow \tau\tau$  validation region is designed in the  $N_{\text{jet}} = 0$  channel to check the normalization and shape of the  $Z \rightarrow \tau\tau$  background in the  $N_{\text{jet}} = 0$  low- $V_{0.3|1.0}$  signal region. This is achieved via the requirement of large opening angle of the charged leptons  $\Delta\Phi(\ell, \ell) > 2.4$ , while at the same time requiring the invariant mass to be below 80 GeV.

In the  $N_{\text{jet}} = 1$  region, a similar invariant mass selection is made, with the addition of a  $m_{\tau\tau} > 66$  GeV selection in order to enrich this region in  $Z \rightarrow \tau\tau$ . This is designed as a control region and used to control the normalization of the  $Z \rightarrow \tau\tau$  background in the  $N_{\text{jet}} = 1$  low- $V_{0.3|1.0}$  regions.



**Figure 10.11:** The  $m_{\ell\ell}$  and  $m_{jj}$  distributions in the  $N_{\text{jet}} = 0$ ,  $N_{\text{jet}} \geq 2$  top control regions, respectively. For the  $N_{\text{jet}} = 1$  top control region, the  $m_T^\ell$  and leading-lepton  $p_T$  distributions are shown in order to illustrate the lepton modelling. The uncertainty corresponds to the total theoretical, experimental and statistical uncertainty.

Since the  $Z \rightarrow \tau\tau$  background is not a dominant background in the  $N_{\text{jet}} \geq 2$  region, the  $N_{\text{jet}} \geq 2$   $Z \rightarrow \tau\tau$  background is estimated via Monte Carlo. The  $m_T$  and  $N_{\text{jet}} = 1$   $e\mu/\mu e$  DNN distribution are shown in Fig. 10.12 for the  $N_{\text{jet}} = 0$  and  $N_{\text{jet}} = 1$  control regions and show good agreement with the modelling.



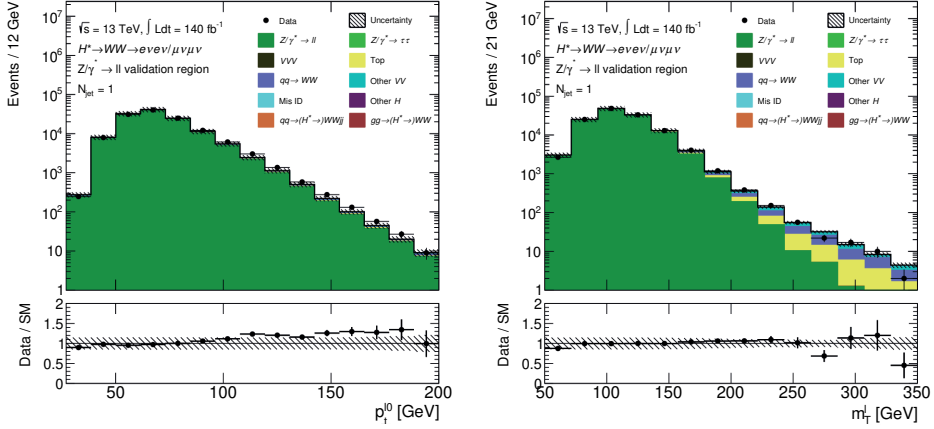
**Figure 10.12:** The  $m_T$  and  $N_{\text{jet}} = 1$   $e\mu/\mu e$  DNN distribution for the  $N_{\text{jet}} = 0$  and  $N_{\text{jet}} = 1$   $Z \rightarrow \tau\tau$  control regions, respectively. The uncertainty corresponds to the total theoretical, experimental and statistical uncertainty.

### 10.7.5 The $Z \rightarrow \ell\ell$ validation regions

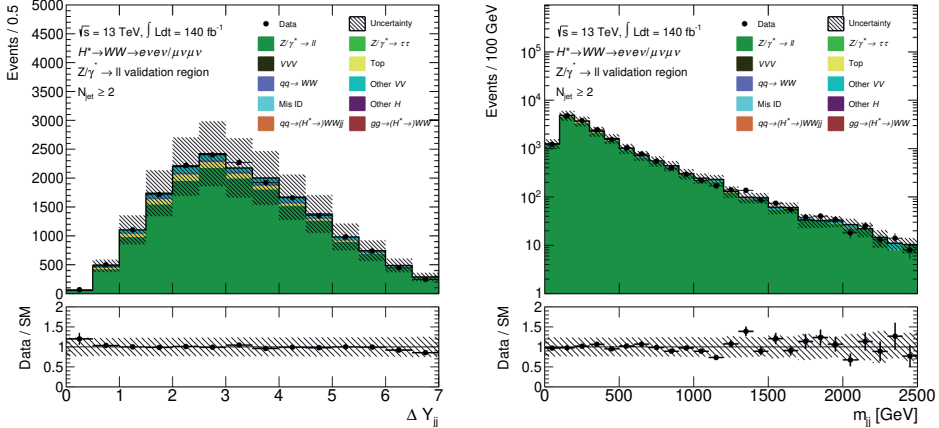
The  $Z \rightarrow \ell\ell$  background is the largest background in the  $ee/\mu\mu$  channels, where it populates the entire  $V_{0.3|1.0}$  distribution. However, it can easily be separated from the signal through the use of the DNN. This results in the fact that no dedicated  $Z \rightarrow \ell\ell$  control region is used in the final fit. The  $Z \rightarrow \ell\ell$  background normalization is estimated via the use of the background-enriched DNN regions.

In order to validate the modelling of this process, a validation region is setup in the  $ee/\mu\mu$  channels. The  $Z \rightarrow \ell\ell$  validation region is defined by the same selection criteria as the signal region, except for a  $|m_{\ell\ell} - 91\text{GeV}| < 15\text{ GeV}$  requirement, selecting the events on the Z-peak. The leading-lepton  $p_T$  and  $\max(m_T^W)$  distributions for the  $N_{\text{jet}} = 1$   $Z \rightarrow \ell\ell$  validation region are shown in Fig. 10.13.

A mis-modelling is present in the leading-lepton  $p_T$  distribution and is caused by a bug in the shower of the SHERPA Monte Carlo simulation. In order to remedy this, a dedicated uncertainty is applied to the  $Z \rightarrow \ell\ell$  background in the fit, which allows the fit to correct for this mis-modelling. More on this method can be found in the systematic uncertainties section. The  $\Delta Y_{jj}$  and  $m_{jj}$  distributions for the  $N_{\text{jet}} \geq 2$   $Z \rightarrow \ell\ell$  validation region are shown in Fig. 10.14, which show good agreement with the data.



**Figure 10.13:** The lepton  $p_T$  and  $\max(m_T^W)$  distributions for the  $N_{\text{jet}} = 1$   $Z \rightarrow \ell\ell$  validation region. The uncertainty corresponds to the total theoretical, experimental and statistical uncertainty.



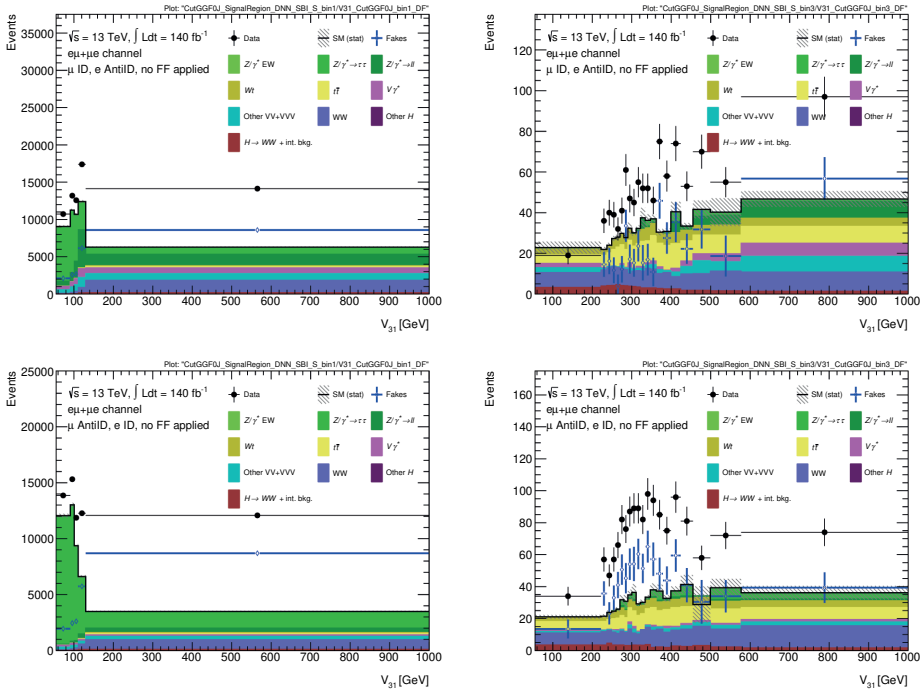
**Figure 10.14:** The  $\Delta Y_{jj}$  and  $m_{jj}$  distributions for the  $N_{\text{jet}} \geq 2$   $Z \rightarrow \ell\ell$  region. The uncertainty corresponds to the total theoretical, experimental and statistical uncertainty.

### 10.7.6 Anti-ID control regions

The  $W$ +jets background contribution is estimated with the data-driven method described in § 9.3.8. Kinematic distributions of the control samples for the  $V_{0.3|1.0}$  distribution in the  $N_{\text{jet}} = 0$   $e\mu/\mu e$  region are shown in Fig. 10.15 and require either an electron or muon fake. Since no fake factors are applied in this figure, the difference between the data and the MC stack corresponds to the fakes yield taken into account by the data-driven method.

From these anti-ID regions, defined using the lepton definitions in Tab. 9.3, the single and double fakes yield are extrapolated into the  $N_{\text{jet}} = 0$   $e\mu/\mu e$  signal region. All regions in the analysis have their own dedicated anti-ID region used for the fake lepton extrapolation.



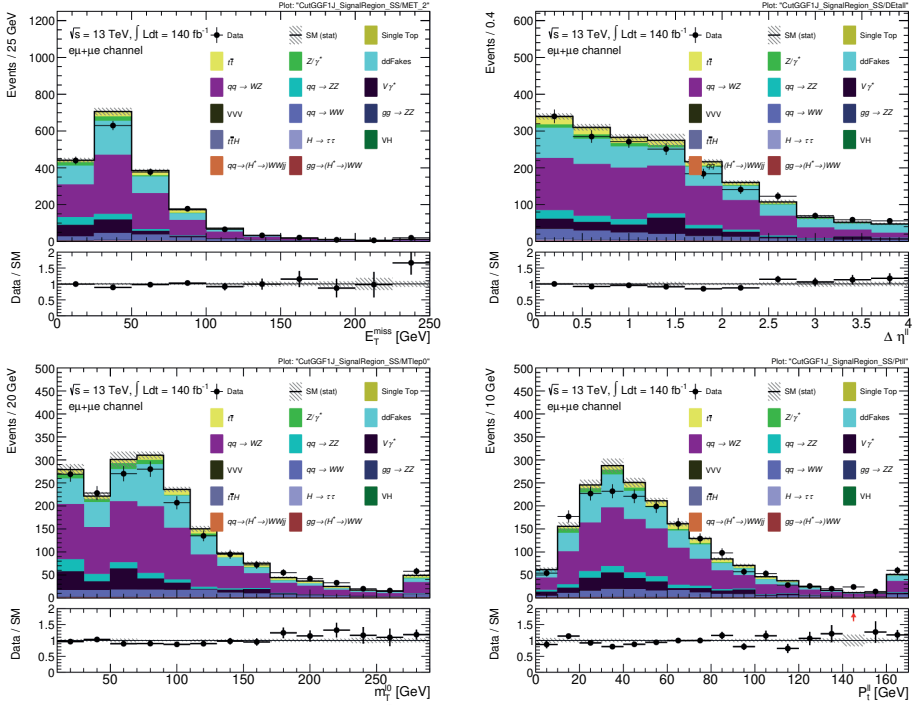


**Figure 10.15:** The fake muon and fake electron  $V_{0.3|1.0}$  distribution in the  $W$ +jets  $N_{\text{jet}} = 0$   $e\mu/\mu e$  control region for the background and signal-enriched DNN regions (without the application of the fake factors). Therefore, the difference between the data and the MC stack corresponds to the fakes yield taken into account by the data-driven method. The unstacked blue histogram labelled fakes corresponds to the difference between the MC stack and the data, and therefore the amount of fakes in this region. The uncertainty corresponds to the statistical uncertainty.

### 10.7.7 Same-sign validation regions

The signal and control regions require the charges of the lepton pairs to be opposite sign, because the SM Higgs boson is neutral. The same-sign validation region is selected like the signal region, except that same instead of opposite lepton charges are required. Requiring same-sign leptons ( $--$  or  $++$ ) rejects the signal and the  $WW$ , top, and  $Z/\gamma^* + \text{jets}$  backgrounds, leaving  $V\gamma$ ,  $W\gamma^*$ ,  $WZ$ , and  $ZZ$ , along with events that involve fake leptons.

The same-sign validation regions offer a good opportunity to study the modelling for these processes, which are small but non-negligible for the SR. The most important kinematic variables related to the modelling of leptons in the  $N_{\text{jet}} = 1$   $e\mu/\mu e$  signal region are shown in Fig. 10.16. The modelling is adequate, considering systematic uncertainties of order 30% are typically applied on the fakes estimation. Notably, there are no trends visible in the  $\eta$ -related distributions, which indicates that the use of  $\eta$ -inclusive fake factors is appropriate.



**Figure 10.16:**  $ggF N_{\text{jet}} = 1 e\mu/\mu e$  same-sign validation region. Hashed bands indicate the statistical uncertainties of the prediction (without the statistical uncertainties of the NFs and fake factors). The uncertainty corresponds to the statistical uncertainty.

## 10.8 Uncertainty estimation

All Monte Carlo samples in this analysis have statistical uncertainties because the number of generated events affects how well the cross-section are estimated. In addition to the statistical uncertainties, the Monte Carlo samples are also subject to systematic uncertainties. The systematic uncertainties are divided into two categories, experimental and theoretical uncertainties. The experimental uncertainties typically originate from the reconstruction and performance of the detector. The theoretical uncertainties originate from assumptions on the modelling of the events.

All systematic uncertainties are propagated through the analysis in order to evaluate the effect of these parameters on the observables ( $V_{0.3|1.0}$  and the scale factors used in the control regions). For some uncertainties, only one-sided variation histograms are used, in the absence of a second alternate template, the effect of the uncertainty is evaluated by taking the difference between the nominal and the one-sided variations and symmetrizing this. Additionally, some uncertainties are smoothed out, since sparsely populated regions of phase-space can lead to statistical fluctuations resulting in unwanted effects in the parameterization of the uncertainty. In order to remove potential problematic effects, a smoothing algorithm using a Gaussian kernel is applied, under the assumption that any

uncertainty effect should be monotonic and correlated across bins. An impression of the expected effect of the systematic uncertainties on the off-shell signal strength parameter  $\mu_{\text{off-shell}}$  is shown in Tab. 10.12 in the results section.

### 10.8.1 Experimental uncertainties

The experimental systematic uncertainties are determined by dedicated performance measurements of the detector and the reconstruction of the events. The systematic uncertainties are evaluated by varying the parameters of the detector and the reconstruction of the events within their uncertainties. By evaluating the reconstruction settings at the nominal values and at the varied values corresponding to  $\pm 1\sigma$  of those settings, the effect of the systematic uncertainties are estimated. A summary of the experimental uncertainties is shown in Tab. 10.6. There are cases where the number of experimental uncertainties on the reconstruction of an object is high and highly correlated. In these cases, a principle component analysis is performed in order to obtain uncorrelated orthogonal uncertainties.

#### Luminosity

The integrated luminosity measurement for the full Run-2 data taken period from 2015 to 2018 measures a total integrated luminosity of  $140 \text{ fb}^{-1}$ . The uncertainty in the combined 2015-2018 integrated luminosity is 0.83% [137], obtained using the LUCID-2 detector [222] for the primary luminosity measurements, complemented by measurements using the inner detector and calorimeters.

#### Pile-up

The evaluation of uncertainty on the pile-up  $\langle\mu\rangle$  rescaling, the data scale factor (1.0/1.03) is varied upward (1./0.99) and downward (1./1.07). This uncertainty is propagated into the pile-up jet rejection tool and jet systematics.

#### Lepton reconstruction, identification, and isolation

The uncertainties estimated for the reconstruction, identification, and isolation of electrons and muons are considered as uncorrelated between the different lepton flavours. Scale factors are estimated to correct for the efficiency on data compared to Monte Carlo for electrons and muons [180]. For electrons, the uncertainty on the energy scale and resolution are also considered [179].

Muons carry an uncertainty to cover the track-to-vertex association split, into a systematic and statistical component. On top of this, the uncertainties on the muon momentum scale and resolution are considered [183]. The systematic uncertainty from  $\tau$ -leptons only enter this measurement via the  $E_T^{\text{miss}}$  calculation and are not taken into account because of their negligible impact.

**Table 10.6:** Summary of the experimental systematic uncertainties considered. Each of these uncertainties is implemented as a nuisance parameter in the likelihood model using an up and down variation.

| Systematic uncertainty                    | $N_{\text{NPs}}$ | Short description                                                                                                     |
|-------------------------------------------|------------------|-----------------------------------------------------------------------------------------------------------------------|
| Event                                     |                  |                                                                                                                       |
| Luminosity                                | 1                | uncertainty on total integrated luminosity                                                                            |
| Pileup Reweighting                        | 1                | uncertainty on pileup reweighting                                                                                     |
| Electrons (scale-factors)                 |                  |                                                                                                                       |
| EL_EFF_Trigger_Total_INPCOR_PLUS_UNCOR    | 1                | trigger efficiency uncertainty                                                                                        |
| EL_EFF_Reco_Total_INPCOR_PLUS_UNCOR       | 1                | reconstruction efficiency uncertainty                                                                                 |
| EL_EFF_ID_CorrUncertaintyNP               | 16               | ID efficiency uncertainty                                                                                             |
| EL_EFF_ID_SIMPLIFIED_UncorrUncertaintyNP  | 18               | ID efficiency uncertainty                                                                                             |
| EL_EFF_Iso_Total_INPCOR_PLUS_UNCOR        | 1                | isolation efficiency uncertainty                                                                                      |
| Electrons (four-momenta)                  |                  |                                                                                                                       |
| EG_SCALE_ALL                              | 1                | energy scale uncertainty                                                                                              |
| EG_SCALE_AF2                              | 1                | energy scale uncertainty that vanishes because all samples used nominally are full-sim                                |
| EG_RESOLUTION_ALL                         | 1                | energy resolution uncertainty                                                                                         |
| Muons (scale-factors)                     |                  |                                                                                                                       |
| MUON_EFF_TrigStatUncertainty              | 1                | trigger efficiency uncertainty                                                                                        |
| MUON_EFF_TrigSystUncertainty              | 1                |                                                                                                                       |
| MUON_EFF_RECO_STAT                        | 1                | reconstruction and ID efficiency uncertainty for muons with $p_T > 15$ GeV                                            |
| MUON_EFF_RECO_SYS                         | 1                |                                                                                                                       |
| MUON_ISO_STAT                             | 1                | isolation efficiency uncertainty                                                                                      |
| MUON_ISO_SYS                              | 1                |                                                                                                                       |
| MUON_TTVA_STAT                            | 1                | track-to-vertex association efficiency uncertainty                                                                    |
| MUON_TTVA_SYS                             | 1                |                                                                                                                       |
| Muons (four-momenta)                      |                  |                                                                                                                       |
| MUON_ID                                   | 1                | momentum resolution uncertainty from inner detector                                                                   |
| MUON_MS                                   | 1                | momentum resolution uncertainty from muon system                                                                      |
| MUON_SCALE                                | 1                | momentum scale uncertainty                                                                                            |
| MUON_SAGITTA_RHO                          | 1                | charge dependent momentum scale uncertainty                                                                           |
| MUON_SAGITTA_RESBIAS                      | 1                |                                                                                                                       |
| Jets (energy scale)                       |                  |                                                                                                                       |
| JET_EffectiveNP_Detector                  | 2                | energy scale uncertainty from the in situ analyses splits into 8 components                                           |
| JET_EffectiveNP_Mixed                     | 3                |                                                                                                                       |
| JET_EffectiveNP_Modelling                 | 4                |                                                                                                                       |
| JET_EffectiveNP_Statistical               | 6                |                                                                                                                       |
| JET_EtaIntercalibration_Modelling         | 1                | energy scale uncertainty on eta-intercalibration (modelling)                                                          |
| JET_EtaIntercalibration_TotalStat         | 1                | energy scale uncertainty on eta-intercalibrations (statistics/method)                                                 |
| JET_EtaIntercalibration_NonClosure_highE  | 1                | energy scale uncertainty on eta-intercalibrations (non-closure)                                                       |
| JET_EtaIntercalibration_NonClosure_negEta | 1                |                                                                                                                       |
| JET_EtaIntercalibration_NonClosure_posEta | 1                |                                                                                                                       |
| JET_Pileup_OffsetMu                       | 1                | energy scale uncertainty on pile-up (mu dependent)                                                                    |
| JET_Pileup_OffsetNPV                      | 1                | energy scale uncertainty on pile-up (NPV dependent)                                                                   |
| JET_Pileup_PtTerm                         | 1                | energy scale uncertainty on pile-up (pt term)                                                                         |
| JET_Pileup_RhoTopology                    | 1                | energy scale uncertainty on pile-up (density $\rho$ )                                                                 |
| JET_Flavor_Composition                    | 1                | energy scale uncertainty on flavour composition                                                                       |
| JET_Flavor_Response                       | 1                | energy scale uncertainty on samples' flavour response                                                                 |
| JET_BJES_Response                         | 1                | energy scale uncertainty on b-jets                                                                                    |
| JET_PunchThrough_MC16                     | 1                | energy scale uncertainty for punch-through jets                                                                       |
| JET_SingleParticle_HighPt                 | 1                | energy scale uncertainty from the behaviour of high- $p_T$ jets                                                       |
| Jets (energy resolution)                  |                  |                                                                                                                       |
| JET_JER_DataVsMC_MC16                     | 1                | energy resolution uncertainty, each for both MC and pseudo-data                                                       |
| JET_JER_EffectiveNP                       | 12               |                                                                                                                       |
| Jets (vertex tagging)                     |                  |                                                                                                                       |
| JET_JvtEfficiency                         | 1                | JVT efficiency uncertainty                                                                                            |
| JET_fJvtEfficiency                        | 1                | fJVT efficiency uncertainty                                                                                           |
| Jets (tagging)                            |                  |                                                                                                                       |
| FT_EFF_Eigen_B                            | 3                | $b$ -tagging efficiency uncertainties ("BTAG_MEDIUM"): 3 components for $b$ jets, 3 for $c$ jets and 4 for light jets |
| FT_EFF_Eigen_C                            | 3                |                                                                                                                       |
| FT_EFF_Eigen_L                            | 4                |                                                                                                                       |
| FT_EFF_Eigen_extrapolation                | 1                | $b$ -tagging efficiency uncertainty on the extrapolation to high- $p_T$ jets                                          |
| FT_EFF_Eigen_extrapolation_from_charm     | 1                | $b$ -tagging efficiency uncertainty on tau jets                                                                       |
| MET                                       |                  |                                                                                                                       |
| MET_SoftTrk_ResoPara                      | 1                | track-based soft term related longitudinal resolution uncertainty                                                     |
| MET_SoftTrk_ResoPerp                      | 1                | track-based soft term related transverse resolution uncertainty                                                       |
| MET_SoftTrk_Scale                         | 1                | track-based soft term related longitudinal scale uncertainty                                                          |

## Jets

The uncertainties on jets is split into two categories, the Jet Energy Scale (JES) [223] and the Jet Energy Resolution (JER) [224]. Additionally, the uncertainties on the JVT are considered [191] split into the impact on central and forward jets. The JES uncertainty is split into multiple components, ranging from the in-situ calibration, the detector response, the pile-up, the flavour composition and the behaviour of high- $p_T$  jets. They are applied as a function of jet  $p_T$  and  $|\eta|$ . The JER uncertainty is applied as a smearing of the JER.

## b-tagging

Scale factors used to correct for the performance of the DL1r model on data are calculated in a dedicated tagging calibration measurement [195], based on the flavour and  $p_T$  of the tagged jet. For each flavour,  $b$ -,  $c$ -, and light-jets, the uncertainties on the scale factors are calculated separately, using the methods described in Ref. [192]. As in this analysis the  $b$ -tagging is used to identify the  $b$ -jets, these uncertainties are the most impactful. The total uncertainty on the  $b$ -tagging scale factors amounts to about 2% for the  $b$ -tagged regions and 5% for the  $b$ -veto regions.

## Fake leptons

Systematic uncertainties in the data-driven fake lepton estimation include statistical uncertainties from data and Monte Carlo, as well as systematic uncertainties covering electroweak subtraction and flavour composition. These are detailed in § 9.3.8, Tab. 9.6 for di-jets fakes, and Tab. 9.4 for  $Z/\gamma^* + \text{jets}$  fakes. Tab. 10.7 summarizes the considered uncertainties.

**Table 10.7:** Summary of the systematic uncertainties considered for data-driven fakes. NP refers to related nuisance parameters in the statistical analysis.

| Systematic uncertainty                                       | Short description                              |
|--------------------------------------------------------------|------------------------------------------------|
| Primary fake factors (from $Z + \text{jets}$ )               |                                                |
| FF stat uncertainty                                          | 1 NP per FF bin, i.e. 4 per $e/\mu$            |
| Electroweak subtraction                                      | 1 NP per $e/\mu$                               |
| Sample composition (for CF)                                  | 1 NP per $e/\mu$                               |
| Bias correction for fakes firing SLT (from di-jets)          |                                                |
| Stat uncertainty (for 2015)                                  | 1 NP per FF bin, i.e. 6 for $e$ , 8 for $\mu$  |
| Stat uncertainty (for 2016-2018)                             | 1 NP per FF bin, i.e. 14 for $e$ , 8 for $\mu$ |
| Electroweak subtraction                                      | 1 NP per $e/\mu$                               |
| Sample composition (for extrapolation to $V + \text{jets}$ ) | 40%, 1 NP per $e/\mu$                          |
| Bias correction for electron fakes firing DLT (from MC)      |                                                |
| Total uncertainty                                            | 1 NP                                           |

### Missing transverse energy

All uncertainties related to the objects which go into the  $E_T^{\text{miss}}$  calculation are considered when building the  $E_T^{\text{miss}}$ . The soft term  $E_T^{\text{miss}}$  that originates from this calculation is also subject to dedicated uncertainties [196] pertaining to its scale and resolution.

#### 10.8.2 Theoretical uncertainties

Similar to the experimental uncertainties, the theoretical uncertainties reflect the uncertainty in the nominal modelling of the events. All Monte Carlo samples used in this analysis carry a set of theoretical uncertainties which allows them to be varied away from the nominal value in the analysis.

Since the analysis aim is to combine with the on-shell HWW analysis, the analysis benefits from a similar treatment of the theory uncertainties where possible. The theory uncertainties are related to the choice of matrix element, parton showers, hadronization, the choice of QCD renormalization and factorization scales, the choice of PDFs and any other choice of generator settings reflecting the uncertainty in the assumptions of the Monte Carlo simulation.

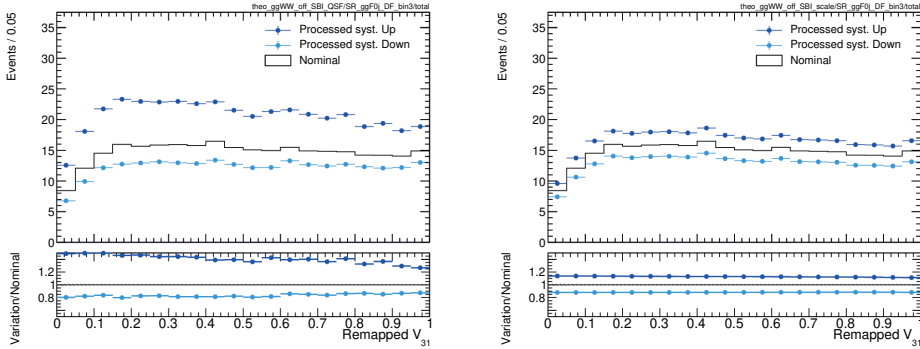
The normalization factors which are derived from control regions are considered modelling uncertainties and are treated as such, since the overall normalization of the Monte Carlo are influenced by the observed data and are therefore considered as a systematic uncertainty. On top of these, there are also *normalization uncertainties* derived fully from the Monte Carlo samples themselves or by comparing to other Monte Carlo samples. Furthermore, the impact of the modelling uncertainties on the shape fit discriminant  $V_{0.3|1.0}$  is evaluated in the form of *shape uncertainties*. Typically, each uncertainty source has a normalization and shape component which are correlated or decorrelated with other uncertainties based on the nature of the uncertainty.

#### Signal modelling uncertainties

For the  $gg \rightarrow (H^*) \rightarrow WW$  and  $qq \rightarrow (H^*) \rightarrow WW + jj$  samples, multiple templates are employed to evaluate the off-shell Higgs signal, including its interference. Consequently, the impact of uncertainties is derived separately for each sample. The SBI, B, and S samples for  $gg \rightarrow (H^*) \rightarrow WW$ , as well as the SBI, SBI20, and B samples for  $qq \rightarrow (H^*) \rightarrow WW + jj$ , are treated individually in the statistical implementation.

***gg-INDUCED WW PRODUCTION:*** The  $gg \rightarrow (H^*) \rightarrow WW$  SHERPA samples are reweighed to their NLO cross-section and the corresponding scale factors are applied as a function of  $m_{WW}^{\text{truth}}$ . The uncertainties on these NLO scale factors are evaluated by comparing the  $m_{WW}/2$  (nominal),  $m_{WW}/4$  (down), and  $m_{WW}$  (up) variations, applied as uncertainties on the nominal prediction. Additionally, the  $\alpha_s$ -scale is varied for the NNPDF3.0nlo set. At particle-level, comparisons are made by comparing the nominal sample with alternative shower variations, which are all SHERPA generator-level settings.

The input variations of the two largest uncertainties, QCD scale and the resummation scale, are shown in Fig. 10.17. The resummation scale uncertainty (QSF) is employed for the resummation of soft gluon emissions, varied by a factor of 0.5 and 2 with respect to the nominal sample. The resummation introduces corrections to the cross-section which are dependent on the ratio between the hard and soft scales of the process. The process has a high hard scale ( $m_{WW}$ ) and a low soft scale ( $p_T^{WW}$ , the boost of the H system) at LO, which is captured by this uncertainty [96]. Additional information on this uncertainty can be found in § 3.4.5.



**Figure 10.17:** The up and down variations of the QSF (left) and scale (right) uncertainties for the  $gg \rightarrow (H^*) \rightarrow WW \rightarrow \ell\nu\ell\nu$  signal template in the  $N_{\text{jet}} = 0$  channel high DNN region.

The matrix element matching scale (CKKW) [225] is used for calculating the overlap between jets from the matrix element and the parton shower. The value in the nominal sample set to 20 GeV. Two samples, with the matching scale set to 15 or 30 GeV, are examined. The default parton shower recoil scheme uses the recoil scheme proposed in [226], while the alternative scheme enables the original recoil strategy documented in [227]. A comparison is made between these two parton showers and used as a two-point uncertainty (CSSKIN). Tab. 10.8 shows the normalization uncertainties for the  $gg \rightarrow (H^*) \rightarrow WW$  samples.

**$qq$ -INDUCED  $WWjj$  PRODUCTION:** The uncertainties on the  $qq \rightarrow (H^*) \rightarrow WW + jj$  SBI, SBI20, and B samples are evaluated separately, but treated as correlated in the profile likelihood fit. For these samples, the  $\mu_R$  and  $\mu_F$  scales are varied using a 7-point variation. In order to account for contributions of channels opening up in higher-order corrections, NLO QCD and EW corrections are calculated for the entire quark-induced  $WWjj$  process, including the off-shell Higgs boson signal, interfering background, and interference [228]. Due to phase-space differences with the fiducial region considered in this calculation, the corrections are not applied directly; rather a 25% uncertainty on the signal, background and interference is used to cover the approximate expected size of the corrections.

**Table 10.8:** Theoretical normalization uncertainties for the ggWW off-shell signal templates in the  $N_{\text{jet}} = 0$  channel.

| Uncertainty Source       | $\pm 1\sigma(\%)$ impact on DNN regions |               |                 |
|--------------------------|-----------------------------------------|---------------|-----------------|
|                          | Background-enriched                     | Mixed         | Signal-enriched |
| S                        |                                         |               |                 |
| CKKW                     | +2.8 / -2.8                             | +3.1 / -3.1   | +3.5 / -3.5     |
| CSSKIN                   | +3.8 / -3.8                             | +7.7 / -7.7   | +5.7 / -5.7     |
| QSF                      | +41.6 / -17.6                           | +36.9 / -16.4 | +32.7 / -13.2   |
| $\alpha_s$               | +1.7 / -1.7                             | +1.7 / -1.7   | +1.6 / -1.6     |
| pdf                      | +1.7 / -1.7                             | +1.9 / -1.9   | +2.0 / -2.0     |
| scale ( $\mu_R, \mu_F$ ) | +18.1 / -14.6                           | +17.5 / -14.4 | +16.5 / -14.3   |
| SBI                      |                                         |               |                 |
| CKKW                     | +7.6 / -7.6                             | +5.8 / -5.8   | +5.0 / -5.0     |
| CSSKIN                   | +4.6 / -4.6                             | +7.8 / -7.8   | +4.7 / -4.7     |
| QSF                      | +51.1 / -20.6                           | +47.4 / -19.0 | +40.4 / -16.7   |
| $\alpha_s$               | +1.5 / -1.5                             | +1.5 / -1.5   | +1.5 / -1.5     |
| pdf                      | +1.8 / -1.8                             | +2.0 / -2.0   | +2.0 / -2.0     |
| scale ( $\mu_R, \mu_F$ ) | +14.3 / -12.3                           | +13.8 / -12.1 | +12.8 / -11.8   |
| B                        |                                         |               |                 |
| CKKW                     | +8.5 / -8.5                             | +6.7 / -6.7   | +6.5 / -6.5     |
| CSSKIN                   | +4.5 / -4.5                             | +7.5 / -7.5   | +4.4 / -4.4     |
| QSF                      | +49.5 / -20.7                           | +45.5 / -18.7 | +37.0 / -15.8   |
| $\alpha_s$               | +1.5 / -1.5                             | +1.5 / -1.5   | +1.5 / -1.5     |
| pdf                      | +1.8 / -1.8                             | +1.9 / -1.9   | +2.0 / -2.0     |
| scale ( $\mu_R, \mu_F$ ) | +14.6 / -12.2                           | +14.0 / -12.1 | +13.2 / -12.1   |

The standard deviation of the NNPDF3.0nlo set is considered and the  $\alpha_s$  scale is varied in the uncertainty derivation. Furthermore, the recommended PYTHIA 8 shower weights are varied. A 7-point variation of the FSR and ISR is applied, varying the FSR and ISR renormalization scales independently. Moreover, a variation of the A14 tuned set of parameters is performed using the Var3cUp and Var3cDown parameter sets. Lastly, a variation of the scale of the non-singular terms is implemented in the FSR and ISR splitting functions. The impact of the normalization uncertainties on the  $q\bar{q} \rightarrow (H^*) \rightarrow WW + jj$  samples is shown in [Tab. 10.9](#).



**Table 10.9:** Theoretical normalization uncertainties for the  $qq \rightarrow (H^*) \rightarrow WW + jj \rightarrow \ell\nu\ell\nu + jj$  off-shell signal templates in the  $N_{\text{jet}} \geq 2$  channel.

| Uncertainty<br>source    | $\pm 1\sigma(\%)$ impact on DNN regions |             |                 |
|--------------------------|-----------------------------------------|-------------|-----------------|
|                          | Background-enriched                     | Mixed       | Signal-enriched |
| SBI20                    |                                         |             |                 |
| Var3c                    | +0.4 / -0.4                             | +1.3 / -1.5 | +2.1 / -2.4     |
| $\alpha_s$               | -0.5 / +0.5                             | -0.3 / +0.3 | -0.1 / +0.1     |
| fsr                      | +2.3 / -2.3                             | +4.3 / -4.3 | +1.7 / -1.7     |
| hardHiLo                 | +0.2 / -0.2                             | +0.8 / -0.7 | +0.2 / -0.2     |
| isr                      | +0.4 / -0.4                             | +1.8 / -1.8 | +2.8 / -2.8     |
| pdf                      | +1.3 / -1.3                             | +1.7 / -1.7 | +2.1 / -2.1     |
| scale ( $\mu_R, \mu_F$ ) | +3.3 / -3.4                             | +7.3 / -6.6 | +12.4 / -10.2   |
| SBI                      |                                         |             |                 |
| Var3c                    | +0.7 / -0.5                             | +2.4 / -2.7 | +2.6 / -2.8     |
| $\alpha_s$               | -0.5 / +0.5                             | -0.3 / +0.3 | -0.1 / +0.1     |
| fsr                      | +2.2 / -2.2                             | +5.6 / -5.6 | +3.8 / -3.8     |
| hardHiLo                 | +0.5 / -0.4                             | +1.1 / -1.1 | +0.5 / -0.5     |
| isr                      | +0.8 / -0.8                             | +3.1 / -3.1 | +3.3 / -3.3     |
| pdf                      | +1.3 / -1.3                             | +1.8 / -1.8 | +2.3 / -2.3     |
| scale ( $\mu_R, \mu_F$ ) | +2.9 / -3.1                             | +6.9 / -6.2 | +11.1 / -9.3    |
| B                        |                                         |             |                 |
| Var3c                    | +0.5 / -0.5                             | +1.6 / -1.8 | +1.9 / -2.0     |
| $\alpha_s$               | -0.5 / +0.5                             | -0.3 / +0.3 | -0.1 / +0.1     |
| fsr                      | +2.5 / -2.5                             | +7.5 / -7.5 | +3.0 / -3.0     |
| hardHiLo                 | +0.5 / -0.5                             | +0.8 / -0.9 | +0.4 / -0.4     |
| isr                      | +0.6 / -0.6                             | +2.1 / -2.1 | +2.4 / -2.4     |
| pdf                      | +1.3 / -1.3                             | +1.7 / -1.7 | +2.2 / -2.2     |
| scale ( $\mu_R, \mu_F$ ) | +2.9 / -3.1                             | +7.0 / -6.3 | +11.4 / -9.5    |

## Background modelling uncertainties

**$t\bar{t}$  PRODUCTION:** The uncertainties associated with  $t\bar{t}$  samples are assessed through an approach encompassing on-the-fly weights, sample comparisons, and reweighting techniques. Previous investigations have highlighted disagreement between data and predictions in  $t\bar{t}$  events, notably in the top quark  $p_T$  distribution, when benchmarked against NNLO calculations [206]. Recent advancements include NNLO QCD accurate computations alongside NLO EW corrections [229], resulting in refinements applied via simultaneous 3D reweighting. Parton-level kinematic variations due to QCD scale uncertainties are accounted for, using the up and down envelopes from the 7-point  $\mu_R$  and  $\mu_F$  scheme. These uncertainties are propagated through the  $p_T^{top}$ ,  $p_T^{t\bar{t}}$ , and  $m_{t\bar{t}}$  distributions using a reweighting tool to estimate their impact on the final fit variable.

The choice of PDF, influencing EW corrections, is scrutinized with comparisons made between NNPDF3.0QED [230] and PDF4LHCLUXQED [231]. Further uncertainties stem from the reweighting order inversion across variables and predictions involving different QCD emissions, segmented by their impacts on  $m_{t\bar{t}}$  and  $p_T^{top}$  distributions.

**$Wt$  AND  $t\bar{t}$  PRODUCTION:** Beyond the uncertainties mentioned in the  $t\bar{t}$  section, additional modelling systematics encompass both the  $t\bar{t}$  and  $Wt$  processes. These are comparisons between the Powheg+Pythia and Powheg+Herwig samples, variations of  $\mu_R$  and  $\mu_F$  for the initial and final state radiation, the standard deviation of the NNPDF2.3lo PDF set and comparisons between Powheg+Pythia and a PYTHIA 8 sample with altered  $p_{thard}$  parameter, which is a scale choice that determines the hardness definition of an event. This last uncertainty covers an estimation of the showering phase-space coverage. For  $Wt$  processes, NLO scale uncertainties replace reweighting scale uncertainties, with an additional interference uncertainty derived from comparing the resulting distribution of a diagram removal scheme to the one from an alternative diagram subtraction scheme.

**$qqWW$  PRODUCTION:**  $qqWW$  variations on the nominal SHERPA 2.2.12 sample at reco level involve assessing the 7-point envelope for  $\mu_R$  and  $\mu_F$  scale variations, the standard deviation of the NNPDF3.0nnlo set, and  $\alpha_s$  scale variations. Comparisons at truth level include different showering options within the SHERPA 2.2.2 generator (CKKW, QSF, CSSKIN), alongside evaluating NLO EW contributions via additive versus multiplicative corrections.

**$WZ$  /  $ZZ$  /  $V\gamma$ ,  $VVV$  PRODUCTION:** Similar to  $qqWW$  EW, uncertainties for these samples encompass variations in  $\mu_R$  and  $\mu_F$  scales, NNPDF3.0nnlo PDF set,  $\alpha_s$  scale, and comparisons across different generator options within SHERPA 2.2.2. For the  $VVV$  process, a constant 12% uncertainty is used.

**$Z/\gamma^* + \text{JETS}$  PRODUCTION:** Due to topological differences between  $\tau$  decay and light lepton decay, uncertainties for  $Z \rightarrow \tau\tau$  and  $Z \rightarrow \ell\ell$  are evaluated separately. In the  $e\mu/\mu e$  analysis, dominant  $Z/\gamma^* + \text{jets}$  contributions arise from  $Z \rightarrow \tau\tau$ . Therefore, un-

certainties for  $Z \rightarrow \ell\ell$  samples in the  $e\mu/\mu e$  region are estimated from the same jet multiplicity region in the  $ee/\mu\mu$  channel due to statistical constraints. Conversely, uncertainties in the  $Z \rightarrow \tau\tau$   $ee/\mu\mu$  region are largely influenced by  $Z \rightarrow \ell\ell$ , necessitating evaluation in the  $e\mu/\mu e$  region and subsequent transfer.

Assessments for SHERPA 2.2.11/2.2.14 samples involve reco-level variations in  $\mu_R$  and  $\mu_F$  scales, NNPDF3.0nnlo set,  $\alpha_s$  scale, and comparisons across different generator options (CKKW, QSF) within SHERPA 2.2.11/2.2.14. Additionally, NLO EW contribution uncertainties are evaluated by comparing the additive and multiplicative NLO EW corrections. a significant mis-modelling observed in  $p_T$  spectra of the  $Z \rightarrow \ell\ell$  and  $Z \rightarrow \tau\tau$  samples, compared against data in dedicated regions, necessitates a dedicated  $p_T$ -modelling uncertainty with SHERPA 2.2.1/2.2.2 samples to quantify differences as uncertainties, particularly in signal and control regions using a two-point comparison.

## 10.9 Statistical model

In this section, the statistical methodology used to derive the results is described. A likelihood function,  $\mathcal{L}$ , is defined using the distribution of the discriminant variable  $V_{0.3|1.0}$ . In each of the signal regions of the jet multiplicity and lepton flavour channels, this variable is predicted via the use of the Monte Carlo samples. The likelihood built from these observables is a product of Poisson functions over the bins of the  $V_{0.3|1.0}$  in the signal regions. On top of this, the total yield in each of the control regions is described with a Poisson distribution; the control regions enter as single-bin counters.

Systematic uncertainties are parameterized as individual nuisance parameters (NPs), named  $\vec{\theta}$ , modelled by a Gaussian function. The technical implementation is done using the HistFactory package [232] in the RooFit [233] and RooStats [234] framework steered by the CAFCore [235] analysis framework.

### 10.9.1 The likelihood function of the off-shell measurement

The likelihood for the interpretation of the off-shell signal strength measurement is based on the following likelihood definition

$$\mathcal{L}_b^{\text{off-shell}} \left( N_b \middle| \{ \mu_p \}, \vec{\theta} \right) = \prod_{\text{SRs+CRs}} \prod_{b \in \text{bins}} \text{Pois} \left( N_b \middle| \sum_p \mu_p S_{p,b}(\vec{\theta}) + \sum_p \sqrt{\mu_p} I_{p,b}(\vec{\theta}) + B_b^{\text{int.}}(\vec{\theta}) + B_b^{\text{non. int.}}(\vec{\theta}) \right).$$

The likelihood associated with the observation of  $N_b$  events in each bin,  $b$ , within the signal and control regions is modelled as a Poisson distribution. In the likelihood equation,  $B_b^{\text{int.}}$  is the expected total interfering background yield,  $B_b^{\text{non. int.}}$  is the expected non-interfering background yield and  $S_{p,b}$  is the expected number of signal events for each of the Higgs production processes, which interfere with the interfering background resulting in a yield  $I_{p,b}$ . The parameter  $\mu_p$  represents the signal strength, which is the cross-section times the branching ratio of the production process  $p$  (ggF or VBF) divided

by the SM expectation. All yields depend on a set of nuisance parameters  $\vec{\theta}$ , as both the expected signal and background are subject to uncertainties from the modelling of experimental effects and theoretical predictions.

To improve the quality of the modelling, each control region features a separate background normalization factor for the process targeted by that control region. This approach allows an overall normalization to absorb relative yield changes between the regions and reduces the impact of systematic uncertainties between the signal and control regions. The NFs are constrained via the use of control regions. Systematic uncertainties are encoded by constraint terms, representing auxiliary measurements,  $\tilde{\theta}$ , of a nuisance parameter,  $\theta$ , impacting this analysis. Typically, these measurements are modelled as a unit Gaussian distribution and are included in the likelihood by multiplying the measurement likelihood by the product of the constraint term

$$\mathcal{L}_{\text{auxiliary measurement}}(\theta) = \frac{1}{\sqrt{2\pi}} e^{-\frac{1}{2}\theta^2}.$$

The auxiliary measurement term results in a mean and variance of the NPs  $\theta$  of 0, and 1, respectively. Since the Monte Carlo samples used in this analysis have a limited size, the yield in a bin  $b$  is only known within a statistical uncertainty. In order to account for this, samples contributing to a given bin are propagated into the likelihood using the Barlow and Beeston method [236]. The  $\gamma$ -parameters, encoding the statistical uncertainty on the simulation prediction in each bin, allows to modify, for example, the expected non-interfering background yield in each bin,  $B_b^{\text{non. int.}}(\vec{\theta}) \rightarrow B_b^{\text{non. int.}}(\vec{\theta}, \vec{\gamma}_b)$ , and are constrained by a Poisson distribution

$$\mathcal{L}_b^{\text{statistical}} = \text{Poiss}(m_b | \tau_b \gamma_b).$$

In this equation, parameter  $m_b$  is defined as the size of the MC sample and has a statistical uncertainty. The relative uncertainty of  $m_b$ , defined as  $\sigma_b^{\text{rel.}}$ , is captured in the parameter  $\tau_b$ , defined as  $(1/\sigma_p^{\text{rel.}})^2$  [232], to capture the nominal auxiliary measurement. The  $\gamma_b$  NP is chosen such that its variations of  $\pm 1$  translate to  $\pm 1\sigma$  variations of the statistical uncertainty of  $m_b$ . By multiplying the likelihood with the statistical term, the statistical uncertainty is propagated into the likelihood function. The likelihood function is then defined as

$$\begin{aligned} \mathcal{L}_b^{\text{off-shell}}(N_b | \{\mu_p\}, \vec{\theta}, \vec{\gamma}) &= \prod_{\text{SRs+CRs}} \prod_{b \in \text{bins}} \\ &\text{Poiss} \left( N_b \left| \sum_p \mu_p S_{p,b}(\vec{\theta}, \vec{\gamma}_b) + \sum_p \sqrt{\mu_p} I_{p,b}(\vec{\theta}, \vec{\gamma}_b) + B_b^{\text{int.}}(\vec{\theta}, \vec{\gamma}_b) + B_b^{\text{non. int.}}(\vec{\theta}, \vec{\gamma}_b) \right. \right) \\ &\times \prod_{\theta \in \vec{\theta}} \frac{1}{\sqrt{2\pi}} e^{-\frac{1}{2}\theta^2} \times \prod_{b \in \text{bins}} \text{Poiss}(m_b | \tau_b \gamma_b). \end{aligned}$$

Since this analysis utilizes several sample templates to describe the complete off-shell signal, the final probability model consists of a linear addition of three templates for each

measured off-shell production mode, ggF or VBF. Therefore, in the likelihood function, the terms  $\mu_p S_{p,b}(\vec{\theta}, \vec{\gamma}_b) + \sum_p \sqrt{\mu_p} I_{p,b}(\vec{\theta}, \vec{\gamma}_b) + B_b^{\text{int}}(\vec{\theta}, \vec{\gamma}_b)$  are replaced with

$$S_{\text{ggF},b}^{\text{SBI}}(\mu_p, \vec{\theta}) = (\mu_p - \sqrt{\mu_p}) S_b^S(\vec{\theta}, \vec{\gamma}_b) + \sqrt{\mu_p} S_b^{\text{SBI}}(\vec{\theta}, \vec{\gamma}_b) \\ + (1 - \sqrt{\mu_p}) S_b^B(\vec{\theta}, \vec{\gamma}_b).$$

In this equation,  $S_b^{\text{SBI}}(\vec{\theta}, \vec{\gamma}_b)$ ,  $S_b^S(\vec{\theta}, \vec{\gamma}_b)$ ,  $S_b^B(\vec{\theta}, \vec{\gamma}_b)$  represent the yield, depending on  $\vec{\theta}$ , per bin of the SBI, S and B samples, respectively. This formulation is used such that the different parts of the signal parameterization remain statistically independent. For the VBF,  $qq \rightarrow (H^*) \rightarrow WW + jj$ , process, a similar but slightly different parameterization is used because of the lack of a VBF signal sample. The yield of the  $qq \rightarrow (H^*) \rightarrow WW + jj$  (SBI) process per bin, using the samples  $S_b^{\text{SBI}}(\vec{\theta}, \vec{\gamma}_b)$ ,  $S_b^{\text{SBI}20}(\vec{\theta}, \vec{\gamma}_b)$ ,  $S_b^B(\vec{\theta}, \vec{\gamma}_b)$ , is described as

$$S_{\text{VBF},b}^{\text{SBI}}(\mu_p, \vec{\theta}) = \frac{20 \cdot \sqrt{\mu_p} - \sqrt{20} \cdot \mu_p}{20 - \sqrt{20}} S_b^{\text{SBI}}(\vec{\theta}, \vec{\gamma}_b) + \frac{\mu_p - \sqrt{\mu_p}}{20 - \sqrt{20}} S_b^{\text{SBI}20}(\vec{\theta}, \vec{\gamma}_b) \\ + \frac{(\sqrt{\mu_p} - 1)(\sqrt{\mu_p} - \sqrt{20})}{\sqrt{20}} S_b^B(\vec{\theta}, \vec{\gamma}_b).$$

### 10.9.2 Minimization and hypothesis testing

The most probable signal estimate is found by numerically minimizing  $-\ln(\mathcal{L})$  with respect to the signal strength parameters  $\mu_p$ , the nuisance parameters  $\vec{\theta}$ ,  $\vec{\gamma}$ , and the background normalizations, using the MINUIT package [237]. This minimization yields the best-fit value and allows for the calculation of the profile likelihood ratio, which is used to evaluate the consistency of the data with the background-only hypothesis. The profile likelihood ratio, also known as the Neyman-Pearson ratio, is defined as

$$\lambda(\mu) = \frac{\mathcal{L}(\mu, \hat{\vec{\theta}}, \hat{\vec{\gamma}})}{\mathcal{L}(\hat{\mu}, \hat{\vec{\theta}}, \hat{\vec{\gamma}})},$$

where  $\hat{\mu}$ ,  $\hat{\vec{\theta}}$ ,  $\hat{\vec{\gamma}}$  are the parameter values that maximize the likelihood for a given  $\mu$  (the global minimum), and  $\hat{\vec{\theta}}$ ,  $\hat{\vec{\gamma}}$  are the values that conditionally maximize the likelihood under the chosen value of  $\mu$  ( $\mu = 0$ ). Two test statistics [238],  $q_0$  and  $t_\mu$ , are derived from  $\lambda(\mu)$  and are defined as

$$q_0 = \begin{cases} -2 \ln \lambda(0) & \text{if } \hat{\mu} \geq 0 \\ 0 & \text{if } \hat{\mu} < 0 \end{cases} \quad \text{and } t_\mu = 2 \ln \lambda(\mu).$$

The difference between  $q_0$  and  $t_\mu$  is that  $q_0$  considers only upward fluctuations of  $\hat{\mu}$  with respect to  $\mu = 0$  to count against the hypothesis of  $\mu = 0$ , whereas  $t_\mu$  counts fluctuations in either direction against the hypothesis  $\mu$ . The test statistic is thus specifically designed for discovery sensitivity, as negative fluctuations are interpreted as fluctuations of the background, rather than that of the signal.

### Calculation of $p$ -values for discovery

Via the use of  $q_0$ , a  $p$ -value can be estimated which quantifies to the probability to observe a statistical fluctuation of the null-hypothesis that is more extreme than the one observed in the data. This  $p$ -value is given by

$$p_0 = \int_{q_{0,\text{observed}}}^{\infty} f(q_0|\mu = 0) dq_0,$$

where  $f(q_0|\mu = 0)$  is the probability distribution function of the test statistic under the null hypothesis ( $\mu = 0$ ). In the results, when not otherwise specified, the  $p$ -value is also expressed in terms of a significance,  $Z$ , which is calculated as

$$Z = \Phi^{-1}(1 - p),$$

with  $\Phi^{-1}$  the inverse of the cumulative distribution function of a standard Gaussian. It is defined such that a Gaussian-distributed variable that is  $Z$  standard deviations above its mean has an upper-tail probability equal to the  $p$ -value.

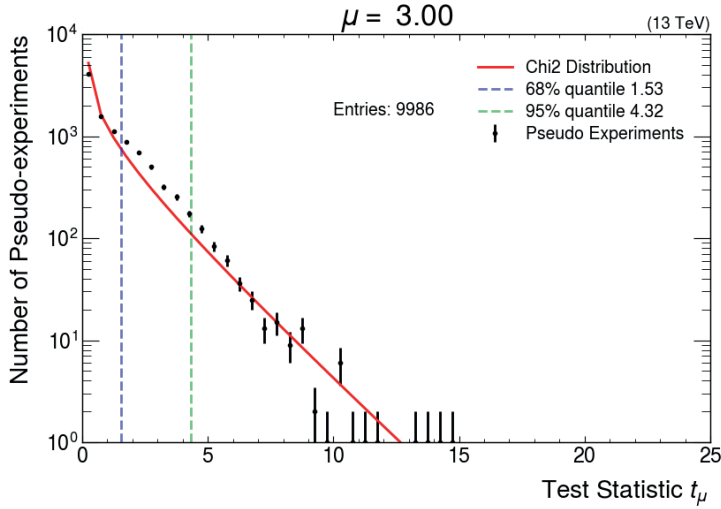
### Construction of confidence intervals

When the best-fit values are estimated by the minimization of the likelihood function, the uncertainties on the POI are derived by profiling the likelihood function and comparing it against the confidence belt of the measurement using the test statistic  $t_\mu$ .

In the asymptotic regime of high statistics, and under some regularity conditions, the distribution of the test statistic  $f(t_\mu|\mu)$  is identical for each value of  $\mu$ , resulting in a rectangular confidence belt for the profile likelihood statistic: a confidence interval can be constructed from a function likelihood ratio  $t(\mu)$  by intersecting it with a fixed threshold of  $t = 0.5 \cdot Z^2$ , where  $Z$  corresponds to the target significance of the interval (e.g.  $Z = 1, 2$  correspond to 68% and 95%, respectively) in this regime. However, the structure of the Higgs off-shell cross-section with interference induces ambiguities in the solutions of the POI, due to the presence of both linear and square-root terms in the parameter of inference, which spoil the regularity conditions. As a result, the distribution of  $f(t_\mu|\mu)$  is no longer  $\chi^2$  distributed, and the shape of the confidence belt is no longer a simple threshold, and must instead be explicitly constructed from simulated data.

To overcome this obstacle inherent to the off-shell analysis, the frequentist procedure of the Neyman construction [239] is followed to construct 1D confidence belts for the off-shell signal strength ( $\mu_{\text{off-shell}}$ ), the total Higgs width ( $\Gamma_H$ ) and the off-shell/on-shell  $VV$  coupling ratio ( $R_{VV}$ ). Any deviation from the asymptotic regime is expected to be negligible for the remaining reported fit results in this thesis. On top of this, the Neyman construction procedure requires a significant computational effort. Therefore, only three 1D confidence belts are constructed.

The procedure of the Neyman construction is as follows. First, multiple toy datasets are created using the statistical model of the analysis. A single toy is constructed from

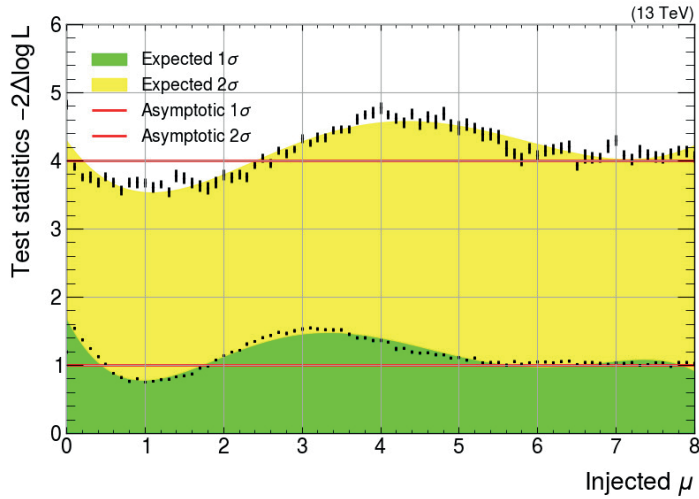


**Figure 10.18:** The test statistic distribution of  $t_\mu$  derived from 9986 toys for an injected  $\mu_{\text{off-shell}}$  signal strength of 3.0.

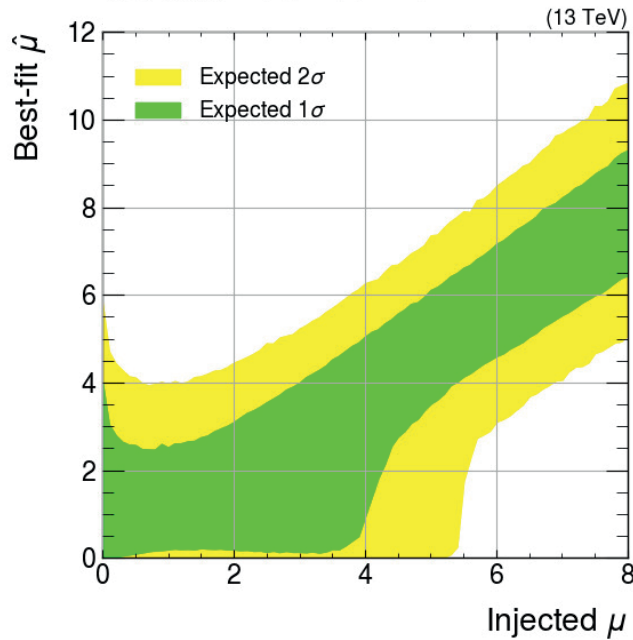
the statistical model with a set, thus known, value of  $\mu$ . Next, the global observables of this toy statistical model are randomized within their expected values. A single toy dataset is then treated as if it were an actual reconstructed dataset and fitted to obtain the measured POI. When the fit on this toy is performed, the observed value for  $\mu$ , denoted as  $\hat{\mu}$ , is found. To construct the confidence belt, the test statistic value  $t_\mu$  is calculated, for 10,000 toy datasets per value of  $\mu$ , in order to numerically obtain the distribution of  $t_\mu$  for a wide range of  $\mu$  values. Fig. 10.18 shows, as example, the distribution of the test statistic  $t_\mu(\mu = 3)$  for an ensemble of datasets sampled under the hypothesis  $\mu = 3$ . The observed distribution of  $t_\mu(\mu = 3)$  deviates significantly from the asymptotically correct  $\chi^2$  distribution visualized by the red curve.

The threshold values for  $\mu = 3$ , for confidence belts at 68% and 95% CL, are obtained by integrating 68% and 95% of the ensemble distribution, respectively, and define the upper bound of the confidence belt at  $\mu = 3.0$ . Fig. 10.19 shows the full confidence belt for  $\mu = [0, 8]$ . The non-Gaussian nature of the measurement is also clearly visible if the confidence belts, constructed with the Neyman construction, are plotted as function  $\hat{\mu}$ , as shown in Fig. 10.20.

Fig. 10.19 also shows that, at values of  $\mu > 5$  and  $\mu > 7$  for 68% and 95% C.L, respectively, the confidence belt becomes asymptotic again. This behaviour is caused by the localised impact of the destructive interference inside the off-shell signal yield causing non-linear effects and shape effects at low values of  $\mu_{\text{off-shell}}$ . The impact of the signal becomes more pronounced at higher values of  $\mu_{\text{off-shell}}$ , as the linear signal-related contributions to the cross-section become more dominant. This, in turn, makes the distribution of the best-fit value of  $\mu_{\text{off-shell}}$  more Gaussian and enter the asymptotic regime.



**Figure 10.19:** The critical values of the test statistic for the off-shell signal strength parameter derived from the Neyman construction procedure compared to the asymptotic approximation. The black lines represent the estimated critical values along with their uncertainty, the yellow and green band borders are determined via a polynomial fit to the data as function of the injected, true, value of  $\mu$ , with which the toys were generated.



**Figure 10.20:** Distribution of the best-fit value of the signal strength modifier  $\hat{\mu}$  as function of the injected, true, value of  $\mu$ , with which the toys were generated.



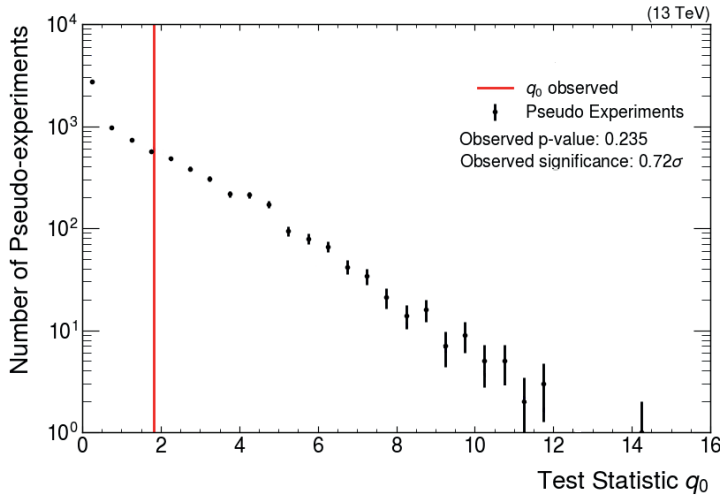
## 10.10 Results

The results of various off-shell signal strength measurements are presented in this section. The results are derived from the likelihood function and minimization techniques described in the previous section. Several fit scenarios are studied, comprised of the 1-POI inclusive off-shell signal strength  $\mu_{\text{off-shell}}$  and the 2-POI  $\mu_{\text{off-shell}}^{\text{ggF}}$  and  $\mu_{\text{off-shell}}^{\text{VBF}}$  combined fit.

The inclusive off-shell signal strength,  $\mu_{\text{off-shell}}$ , is measured in the  $N_{\text{jet}} = 0$ ,  $N_{\text{jet}} = 1$ ,  $N_{\text{jet}} \geq 2$  different and  $ee/\mu\mu$  channels simultaneously. The observed signal strength is found to be

$$\mu_{\text{off-shell}} = 0.27^{+0.53}_{-0.22} (\text{stat.})^{+0.67}_{-0.12} (\text{syst.})^{+0.85}_{-0.25} (\text{tot.}).$$

The observed (expected) signal strength is in agreement with the SM prediction, with a  $p$ -value of 0.24 (0.11) estimated from a distribution of the test statistic  $q_0$  obtained from toy ensemble studies, shown in in Fig. 10.21. A corresponding observed (expected) upper limit on  $\mu_{\text{off-shell}}$  of 3.4 (4.4) is set, at the 95% CL, using the confidence belts built via the Neyman construction<sup>1</sup>.



**Figure 10.21:** The observed significance for rejecting the  $\mu_{\text{off-shell}}$  background-only hypothesis estimated from a distribution of the test statistic  $q_0$  obtained from toy ensemble studies.

It is found that the lower bound of the signal strength is statistically limited, while the higher bound is slightly more limited by the systematic uncertainties. This asymmetric behaviour is expected due to the large theoretical signal uncertainties, which becomes more impactful when the signal strength grows in size and the detectable signal shape

<sup>1</sup> This 95% CL upper limit is derived using the  $t_\mu$  test statistic, which considers fluctuations in both directions. This will overcover when interpreted as a one-sided limit. The quoted limited limit is therefore somewhat conservative.

moves towards the high- $m_{WW}$  proxy regions where the relative systematic uncertainties are significantly larger. The impact of the individual systematic uncertainties is shown in the breakdown plots in § 10.10.1. When the total uncertainty is evaluated, the measurement is statistically limited.

The pre- and post-fit background normalizations of the 1POI fit are listed in Tab. 10.10. With respect to the background only hypothesis the observed (expected) signal strength, or data excess, corresponds with an observed (expected) significance  $Z$  of  $0.7\sigma$  ( $1.2\sigma$ ). The expected significance is evaluated using a profiled Asimov dataset, which is a dataset that is generated using the best-fit values of the nuisance parameters.

**Table 10.10:** Table of observed and expected best-fit normalization factors in the  $\mu_{\text{off-shell}}$  fit.

|          |                             | $qq \rightarrow WW$    | Top-quark              | $Z \rightarrow \tau\tau$ | $Z \rightarrow \ell\ell$ |
|----------|-----------------------------|------------------------|------------------------|--------------------------|--------------------------|
| Expected | $N_{\text{jet}} = 0$ ggF    | $1^{+0.07}_{-0.07}$    | $1^{+0.11}_{-0.10}$    | $1^{+0.06}_{-0.06}$      | $1^{+0.15}_{-0.13}$      |
|          | $N_{\text{jet}} = 1$ ggF    | $1^{+0.15}_{-0.12}$    | $1^{+0.08}_{-0.09}$    | $1^{+0.12}_{-0.10}$      | $1^{+0.21}_{-0.16}$      |
|          | $N_{\text{jet}} \geq 2$ VBF | $1^{+0.6}_{-0.4}$      | $1^{+0.05}_{-0.05}$    |                          | $1^{+0.41}_{-0.27}$      |
| Observed | $N_{\text{jet}} = 0$ ggF    | $1.01^{+0.07}_{-0.07}$ | $0.93^{+0.10}_{-0.09}$ | $0.90^{+0.05}_{-0.05}$   | $1.15^{+0.17}_{-0.15}$   |
|          | $N_{\text{jet}} = 1$ ggF    | $0.90^{+0.14}_{-0.12}$ | $0.97^{+0.08}_{-0.08}$ | $0.90^{+0.11}_{-0.09}$   | $1.08^{+0.23}_{-0.18}$   |
|          | $N_{\text{jet}} \geq 2$ VBF | $0.9^{+0.6}_{-0.4}$    | $0.97^{+0.05}_{-0.05}$ |                          | $0.87^{+0.35}_{-0.23}$   |

The 2-POI fit is performed to measure the off-shell signal strength in the ggF and VBF production modes simultaneously. The observed signal strengths are found to be

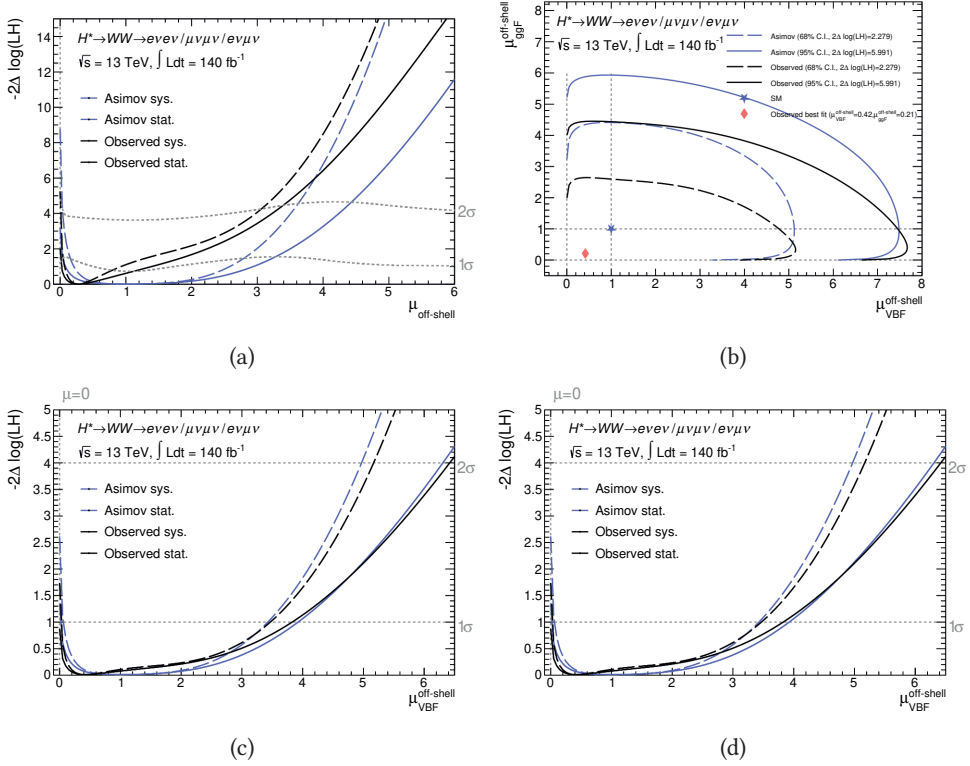
$$\begin{aligned}\mu_{\text{off-shell}}^{\text{ggF}} &= 0.21^{+0.73}_{-0.17} \text{ (stat.)}^{+1.09}_{-0.12} \text{ (syst.)}^{+1.31}_{-0.21} \text{ (tot.)} \\ \mu_{\text{off-shell}}^{\text{VBF}} &= 0.42^{+3.06}_{-0.39} \text{ (stat.)}^{+1.53}_{-0.16} \text{ (syst.)}^{+3.42}_{-0.42} \text{ (tot.)},\end{aligned}$$

where it should be noted that the lower bound of the 68% uncertainty is not reached within parameter space where the model is valid (i.e.  $\mu > 0$ ). For that reason, the quoted lower bound on  $\mu$  is instead the value where the validity boundary is reached. The observed event yield in data, compared to post-fit predictions with  $\mu_{\text{off-shell}}$ , all the nuisance parameters, and normalisation factors set to their best-fit values, is shown in Tab. 10.11.

In order to assess the differential data-to-simulation prediction, the post-fit distributions of the 1-POI fit are shown for all signal regions in Figs. 10.24 to 10.26, showing good agreement between data and simulation. In these plots, the total Monte Carlo simulation is compared to the data. In addition, the unstacked ggF and VBF post-fit SBI yield, corresponding to the best-fit  $\mu_{\text{off-shell}} = 0.27$ , is shown. The likelihood scans made for the 1 and 2-POI fits that produced these results are illustrated in Fig. 10.22.

**Table 10.11:** The observed event yield in data and (total) post-fit predictions, with nuisance parameters and normalization factors set to their best-fit values ( $\hat{\mu}$ ,  $\hat{\theta}$ ), lists of signal and interference ( $S + I$ ), interfering background ( $B$ ), and non-interfering backgrounds for the combined  $e\mu/\mu e$  and  $ee/\mu\mu$  signal regions. The second-to-last row gives the total expected yield, scaled with post-fit nuisance parameters and, for  $S + I$ , the best-fit value of the parameter of interest. The  $1\sigma$  uncertainty includes post-fit systematic and statistical uncertainties combined.

| Process ( $e\mu/\mu e + ee/\mu\mu$ ) | $N_{\text{jet}} = 0$ ggF |                 |               | $N_{\text{jet}} = 1$ ggF |                 |               | $N_{\text{jet}} \geq 2$ VBF |              |              |
|--------------------------------------|--------------------------|-----------------|---------------|--------------------------|-----------------|---------------|-----------------------------|--------------|--------------|
|                                      | low                      | medium          | high          | low                      | medium          | high          | low                         | medium       | high         |
| DNN-score                            |                          |                 |               |                          |                 |               |                             |              |              |
| $H_B$ (ggF)                          | 1747                     | 1828            | 629           | 1052                     | 677             | 190           | 317                         | 21           | 13           |
| $H_{S+I}$ (ggF)                      | -43                      | -107            | -107          | -78                      | -84             | -44           | -39                         | -5           | -3           |
| $H_{EWB}$ (VBF)                      | 13                       | 14              | 9             | 73                       | 66              | 71            | 166                         | 95           | 180          |
| $H_{EW S+I}$ (VBF)                   | 0                        | -2              | 0             | -1                       | -4              | -11           | -8                          | -6           | -17          |
| $t\bar{t}$                           | 4108                     | 6032            | 2566          | 19163                    | 5325            | 1065          | 9316                        | 335          | 142          |
| $q\bar{q} \rightarrow WW$            | 22108                    | 13876           | 2550          | 10393                    | 3811            | 708           | 1879                        | 103          | 46           |
| Mis-Id                               | 2065                     | 772             | 75            | 715                      | 218             | 43            | 202                         | 10           | 7            |
| $Z \rightarrow \tau\tau$             | 65877                    | 692             | 81            | 1446                     | 255             | 20            | 479                         | 6            | 5            |
| $Z \rightarrow \ell\ell$             | 27923                    | 968             | 35            | 9523                     | 335             | 158           | 1057                        | 43           | 7            |
| Other H                              | 600                      | 83              | 3             | 9523                     | 92              | 158           | 1057                        | 43           | 7            |
| $VV(V)$                              | 7424                     | 811             | 124           | 2717                     | 367             | 76            | 451                         | 22           | 6            |
| Total                                | $131823 \pm 992$         | $24967 \pm 217$ | $5966 \pm 58$ | $45096 \pm 740$          | $10986 \pm 113$ | $2277 \pm 33$ | $13869 \pm 467$             | $628 \pm 19$ | $389 \pm 12$ |
| Observed                             | 131851                   | 24983           | 5929          | 44876                    | 11001           | 2281          | 13863                       | 661          | 380          |

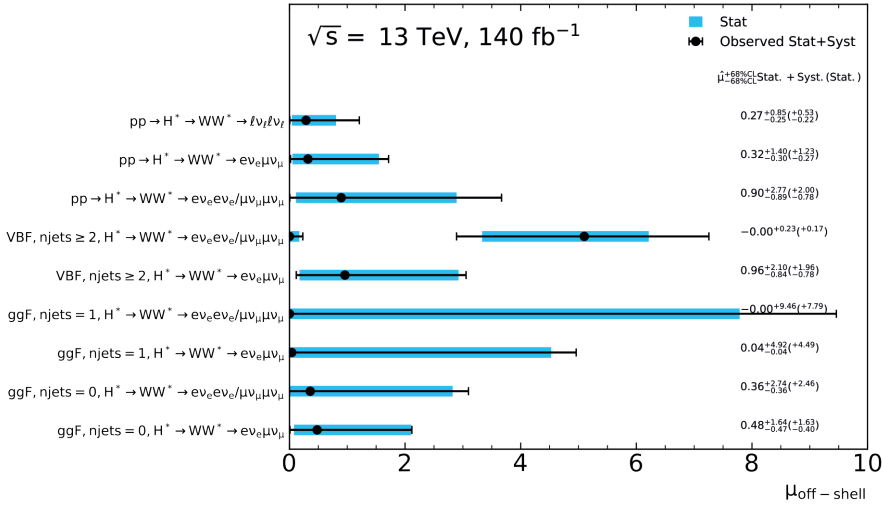


**Figure 10.22:** The signal strength  $\mu_{\text{off-shell}}$  likelihood scan of the 1-POI fit along with the  $1\sigma$  and  $2\sigma$  constructed confidence intervals (a) and the 1D and 2D likelihood scans of the 2-POI fit (b,c,d) using asymptotic  $1\sigma$  and  $2\sigma$  confidence intervals.

A 6-POI fit is performed to measure the off-shell signal strength in the ggF and VBF production modes and their jet categories, simultaneously, and to illustrate the relative measuring strength of each of the 6 jet channels. The observed 6-POI signal strengths are summarized in Fig. 10.23 and compared against the 1-POI fit and a 2-POI fits. It can be seen that between the two flavour channels, the  $e\mu/\mu e$  channel dominates the measurement, mainly due to the  $N_{\text{jet}} = 1$  and the  $N_{\text{jet}} \geq 2$  sub-channels. From the  $ee/\mu\mu$  channel, the  $N_{\text{jet}} = 0$  channel brings the most sensitivity to the total measurement.

The post-fit distributions of the fitted variable  $V_{0.3|1.0}$  are shown for the 0-, 1- and  $\geq 2$ -jet regions in the  $e\mu/\mu e$  and  $ee/\mu\mu$  channels across all DNN regions in Figs. 10.24 to 10.26. The data is in agreement with the Monte Carlo simulation after the fit and the figures show the unstacked SBI yield for both the off-shell ggF and the VBF process for comparison with the total measured yield.

As the measurement uses a finely binned distribution optimized for sensitivity, a more detailed look at the post-fit distributions is given in Figs. 10.27 to 10.29 in terms of the density of events in each bin. This renormalization of the event yields to event density illustrates the lineshape of the  $V_{0.3|1.0}$  distribution in the signal regions, which is of



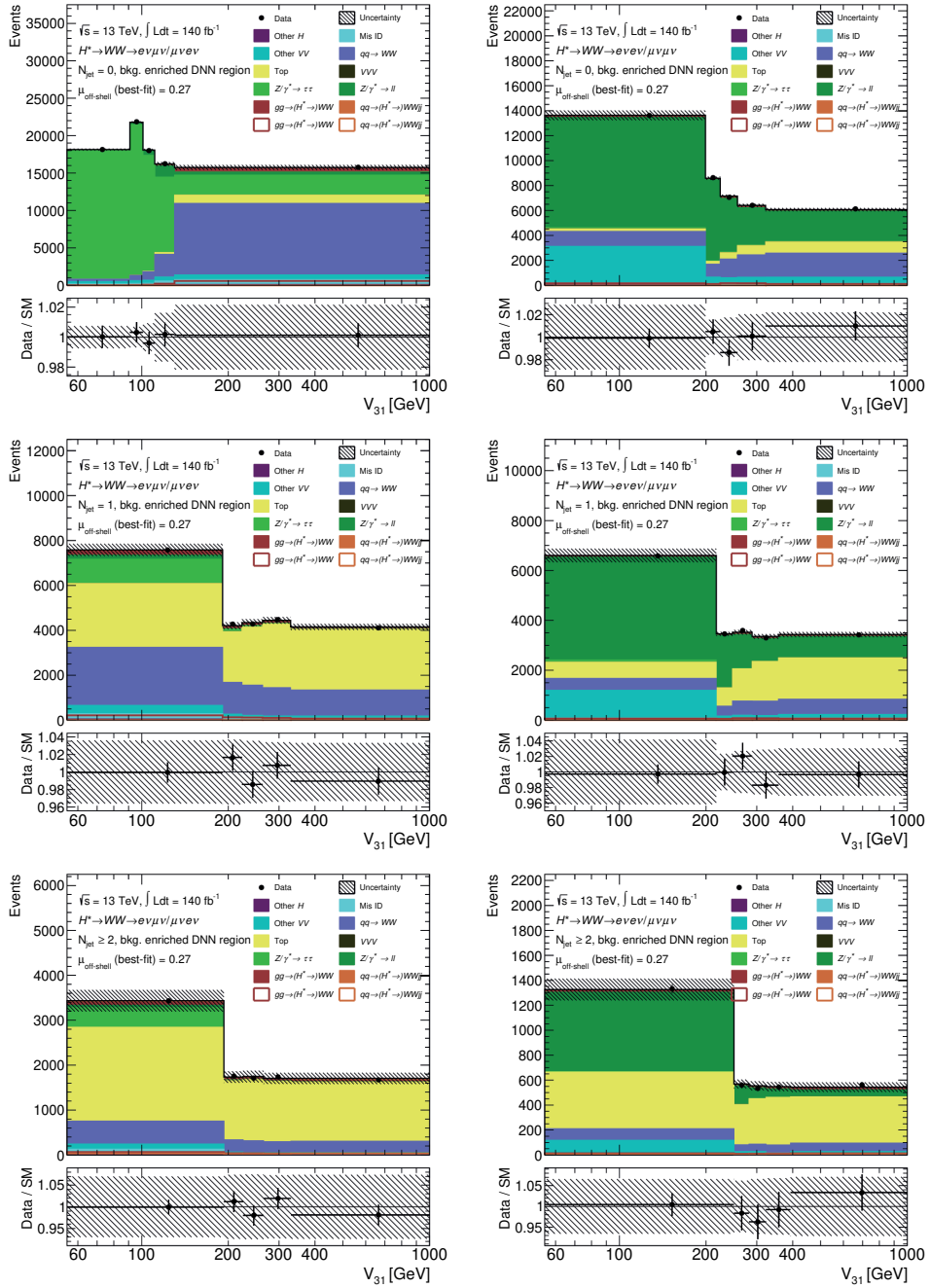
**Figure 10.23:** The observed signal strengths and their observed (expected) uncertainty for the 6-POI fit, the flavour separated and the 1-POI configuration of the analysis along with their total, statistical and systematic uncertainty. In cases where  $\mu_{\text{off-shell}}$  hits the 0 boundary only the statistical uncertainty is shown.

interest in case of off-shell processes.

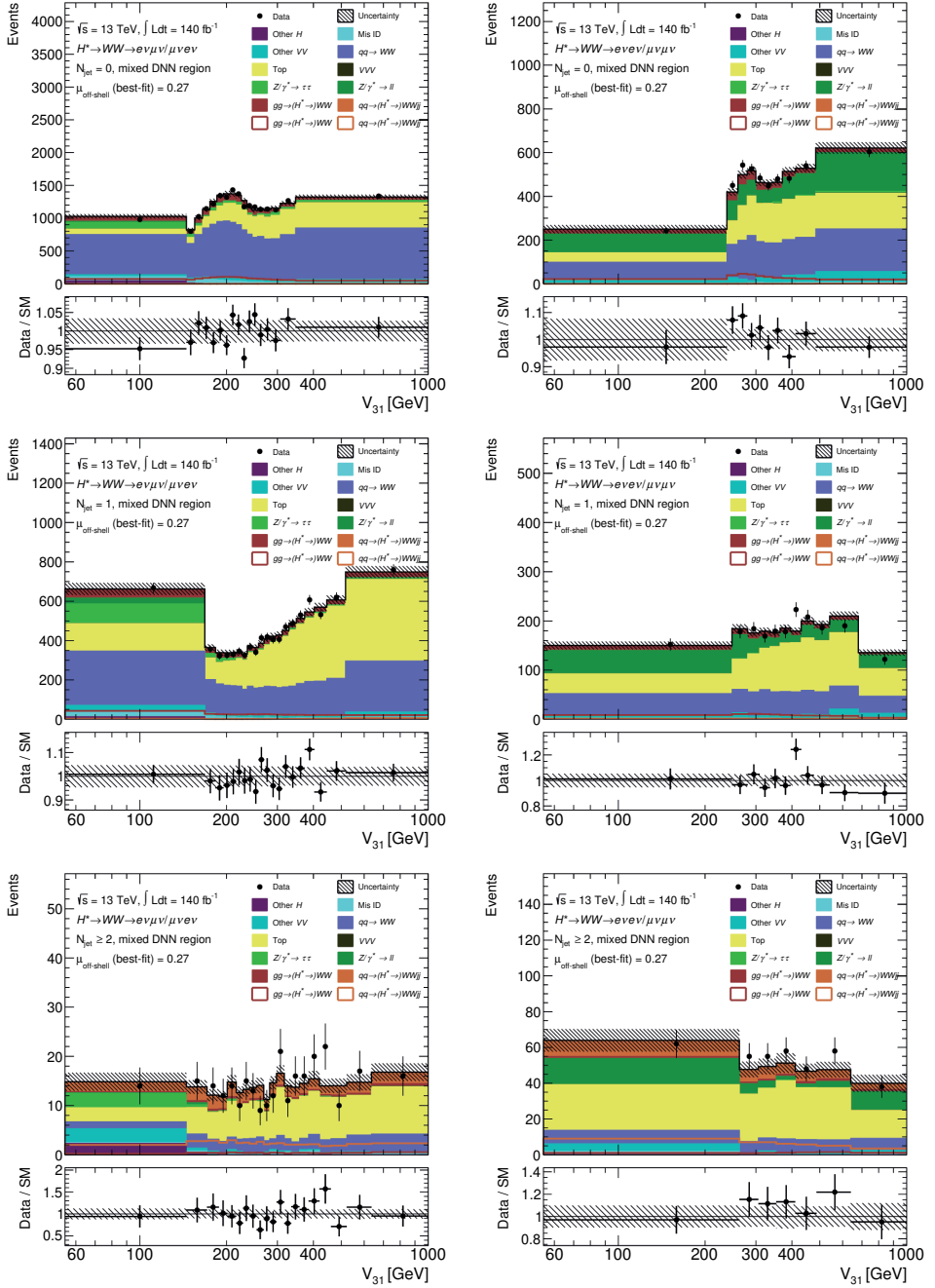
In order to give an indication of the fit quality, goodness of fit metrics such as the  $\chi^2$  per degree of freedom in each of the DNN regions and the corresponding significance of the data-to-simulation agreement in each of the histogram bins is shown. As expected, the  $\chi^2$  per degree of freedom is found to be close to unity in most regions, indicating a good fit. The disagreement significance of the data-to-simulation is highlighted in blue if it exceeds  $1\sigma$  and in red if it exceeds  $2\sigma$ . No data-to-simulation disagreements are found above  $2.7\sigma$  in any of the regions. The number of bins with a disagreement above  $1\sigma$  and  $2\sigma$  agrees with what is expected from statistical fluctuations.

The data excesses, in terms of event density, corresponding to the observed signal strength, are shown in Figs. 10.30 to 10.32 for each unfolded signal region. They are found to be consistent with the expected signal strength within the estimated uncertainties.

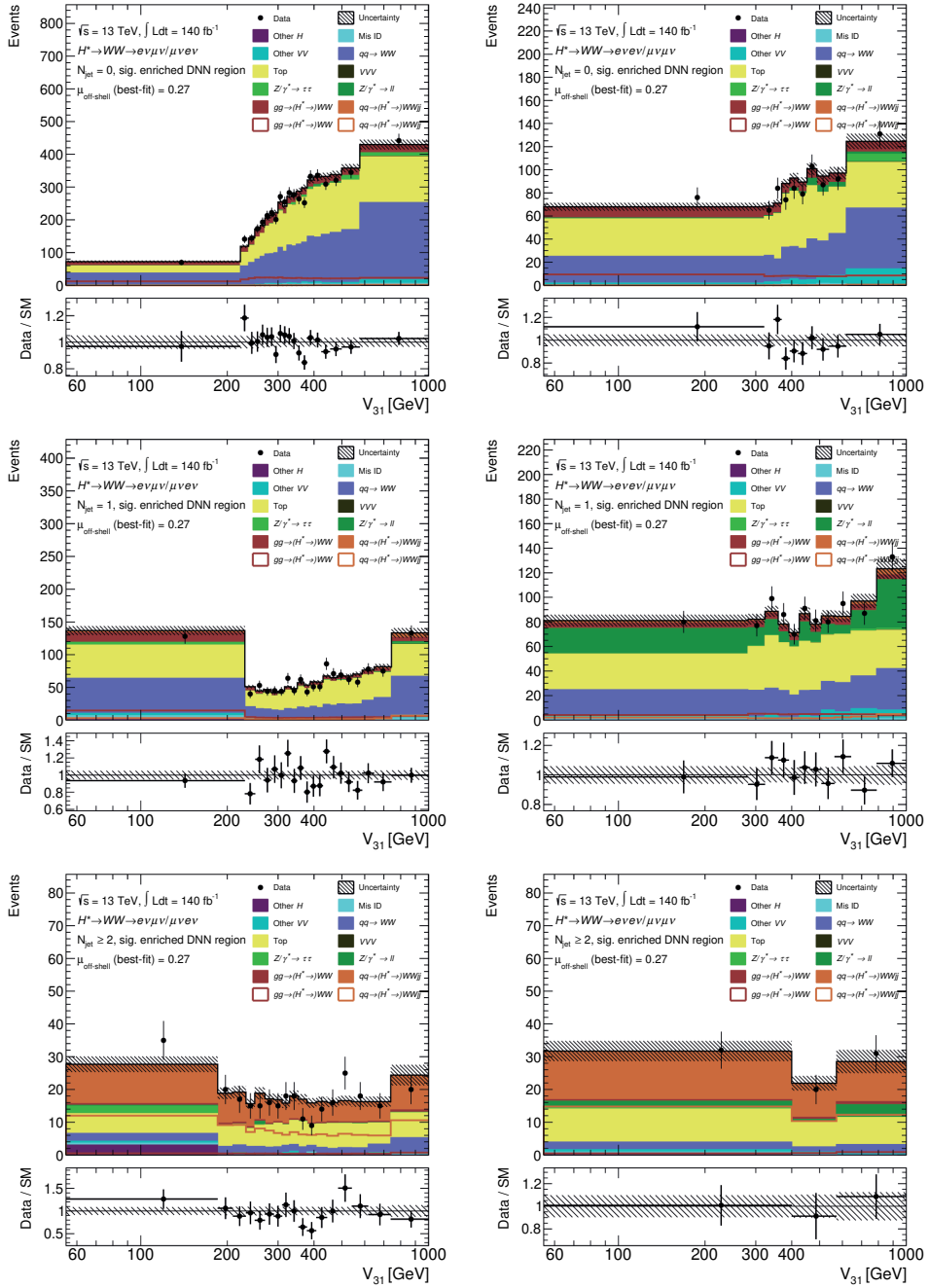
Reconstructed  $N_{\text{jet}} = 0$   $e\mu/\mu e$ ,  $N_{\text{jet}} = 1$   $e\mu/\mu e$ , and  $N_{\text{jet}} \geq 2$   $ee/\mu\mu$  candidate events in each off-shell signal-enriched region appear after Parts II to IV. Electrons and muons are shown as purple and red lines,  $E_{\text{T}}^{\text{miss}}$  as a dotted line, and jets as green cones. The figure displays the electromagnetic barrel calorimeter, with yellow cells for endcap and LAr barrel deposits, green for FCal, and cyan for hadronic endcap deposits.



**Figure 10.24:** The post-fit plots of the  $\mu_{\text{off-shell}}$  signal strength measurement in the background-enriched DNN regions of the  $N_{\text{jet}} = 0$ ,  $N_{\text{jet}} = 1$  and  $N_{\text{jet}} \geq 2$   $e\mu/\mu\mu$  and  $ee/\mu\mu$  channels. The uncertainty in the top and ratio pad corresponds to the total post-fit theoretical, experimental and statistical uncertainty.

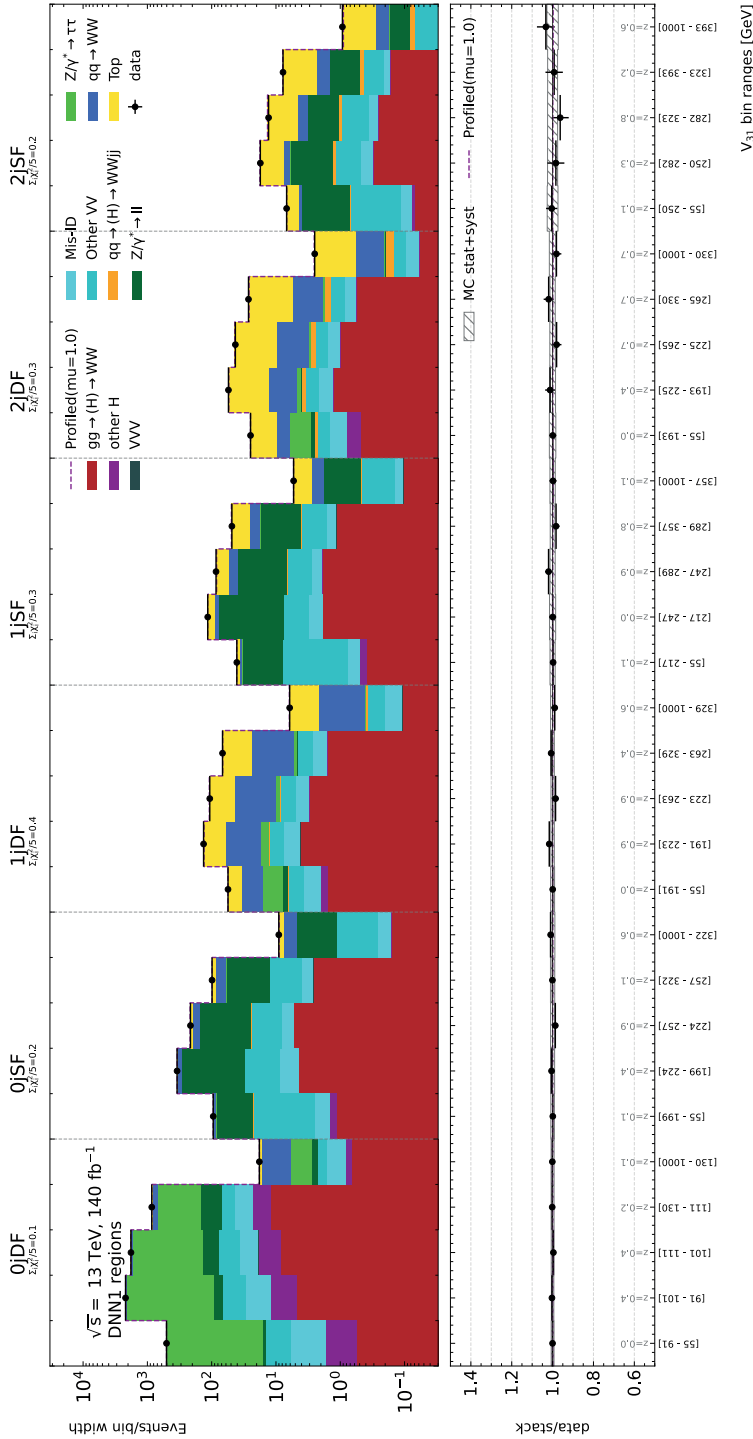


**Figure 10.25:** The post-fit plots of the  $\mu_{\text{off-shell}}$  signal strength measurement in the mixed DNN regions of the  $N_{\text{jet}} = 0$ ,  $N_{\text{jet}} = 1$  and  $N_{\text{jet}} \geq 2$   $e\mu/\mu\mu$  and  $ee/\mu\mu$  channels. The uncertainty in the top and ratio pad corresponds to the total post-fit theoretical, experimental and statistical uncertainty.

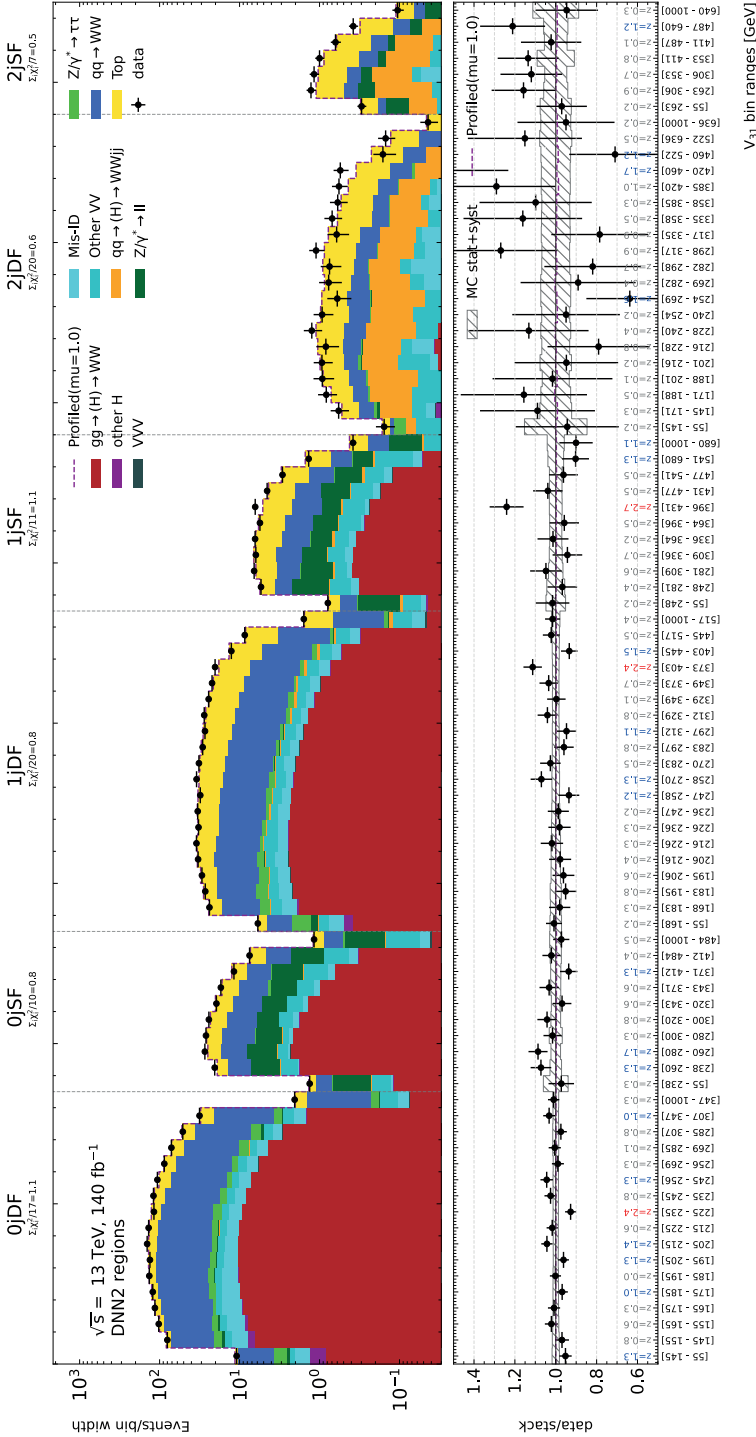


**Figure 10.26:** The post-fit plots of the  $\mu_{\text{off-shell}}$  signal strength measurement in the signal-enriched DNN regions of the  $N_{\text{jet}} = 0$ ,  $N_{\text{jet}} = 1$  and  $N_{\text{jet}} \geq 2$   $e\mu/\mu e$  and  $ee/\mu\mu$  channels. The uncertainty in the top and ratio pad corresponds to the total post-fit theoretical, experimental and statistical uncertainty.

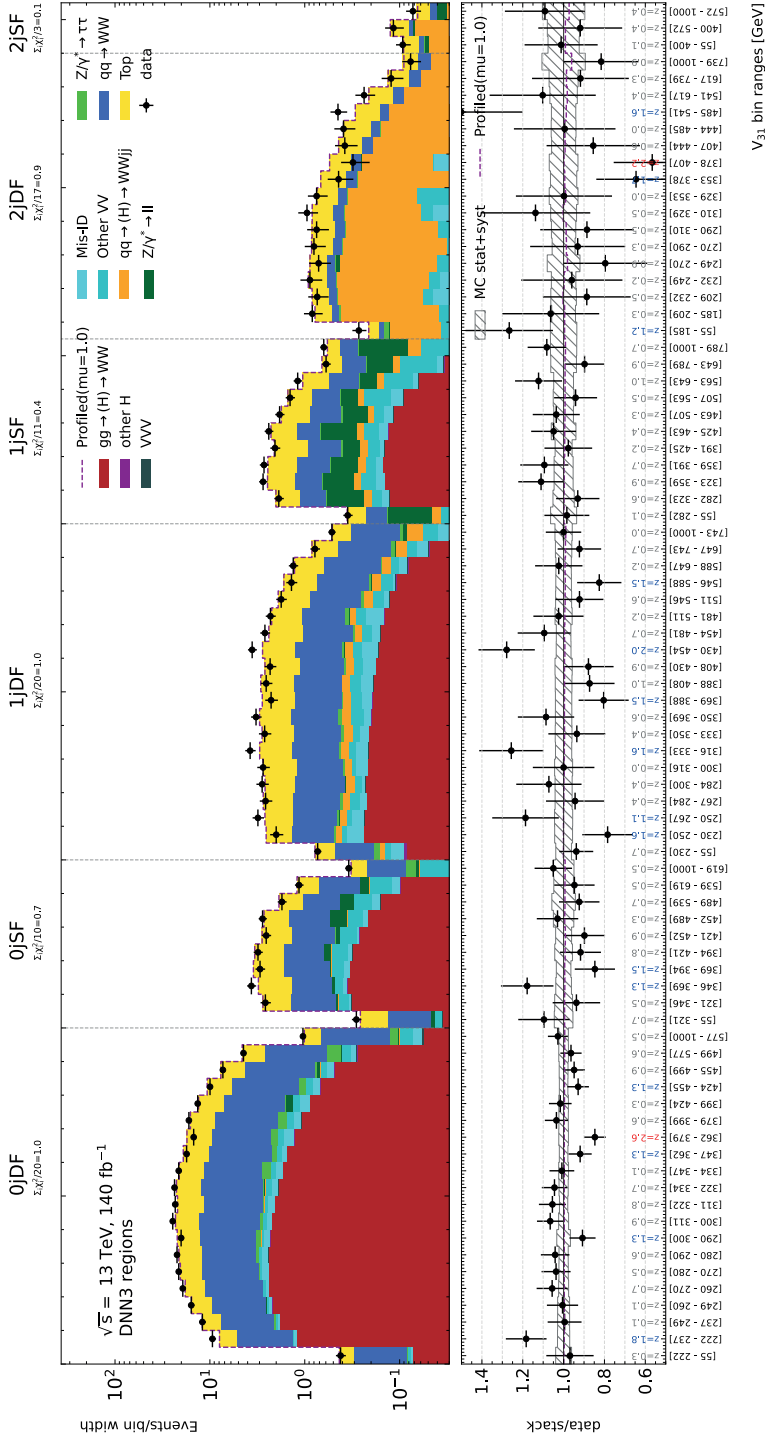




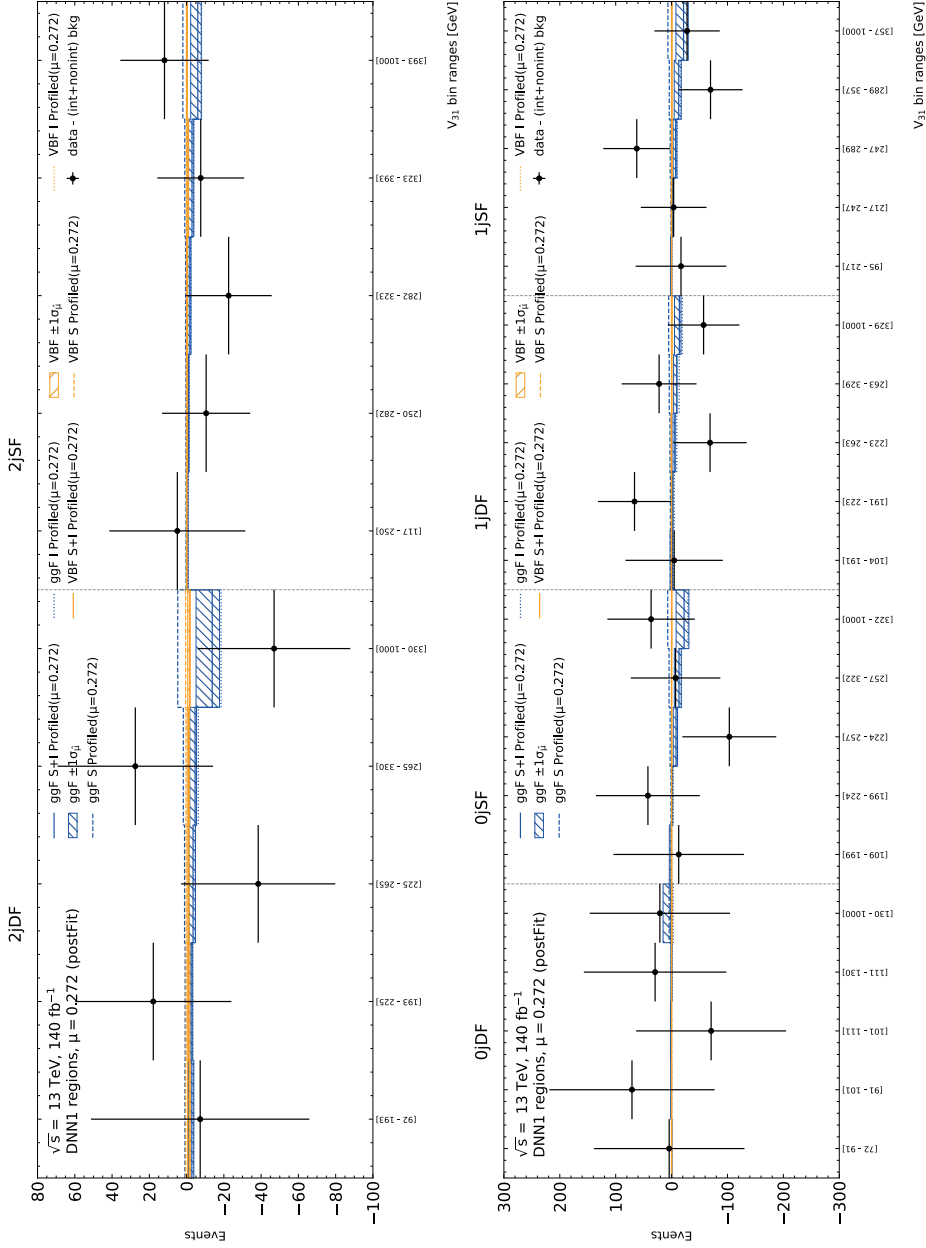
**Figure 10.27:** The post-fit plots of the  $\mu_{\text{off-shell}}$  signal strength measurement in the background-enriched DNN regions. The y-axis is the density of events in each bin and the x-axis is the  $V_{0.3|1.0}$  variable. The goodness of fit metrics are shown in the top of each plot region along with the significance of the data points in each of the bins in grey. The uncertainty in the top and ratio pad corresponds to the total post-fit theoretical, experimental and statistical uncertainty.

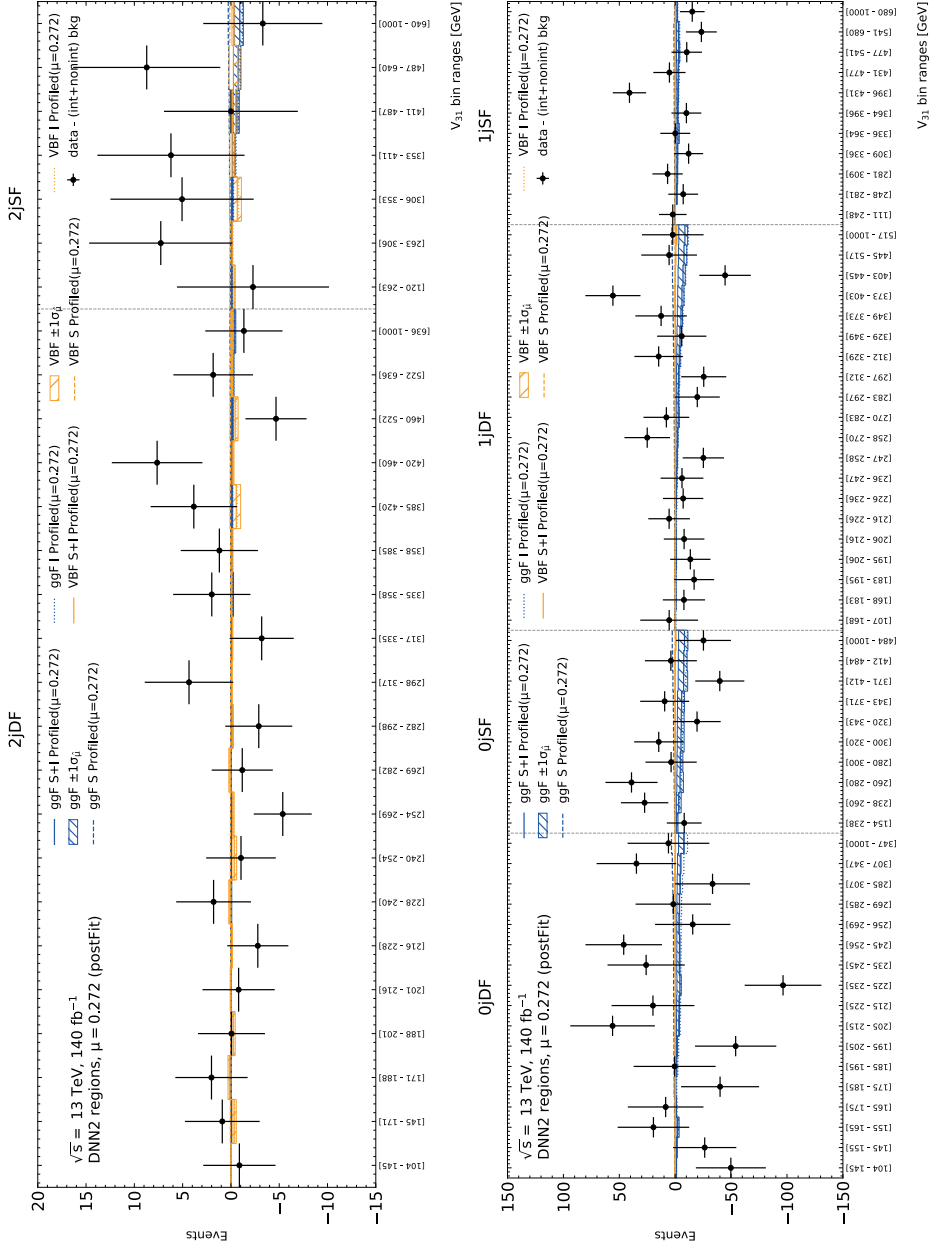


**Figure 10.28:** The post-fit plots of the  $\mu_{\text{off-shell}}$  signal strength measurement in the mixed DNN regions. The y-axis is the density of events in each bin and the x-axis is the  $V_{0.3|1.0}$  variable. The goodness of fit metrics are shown in the top of each plot region along with the significance of the data points in each of the bins in grey. The uncertainty in the top and ratio pad corresponds to the total post-fit theoretical, experimental and statistical uncertainty.

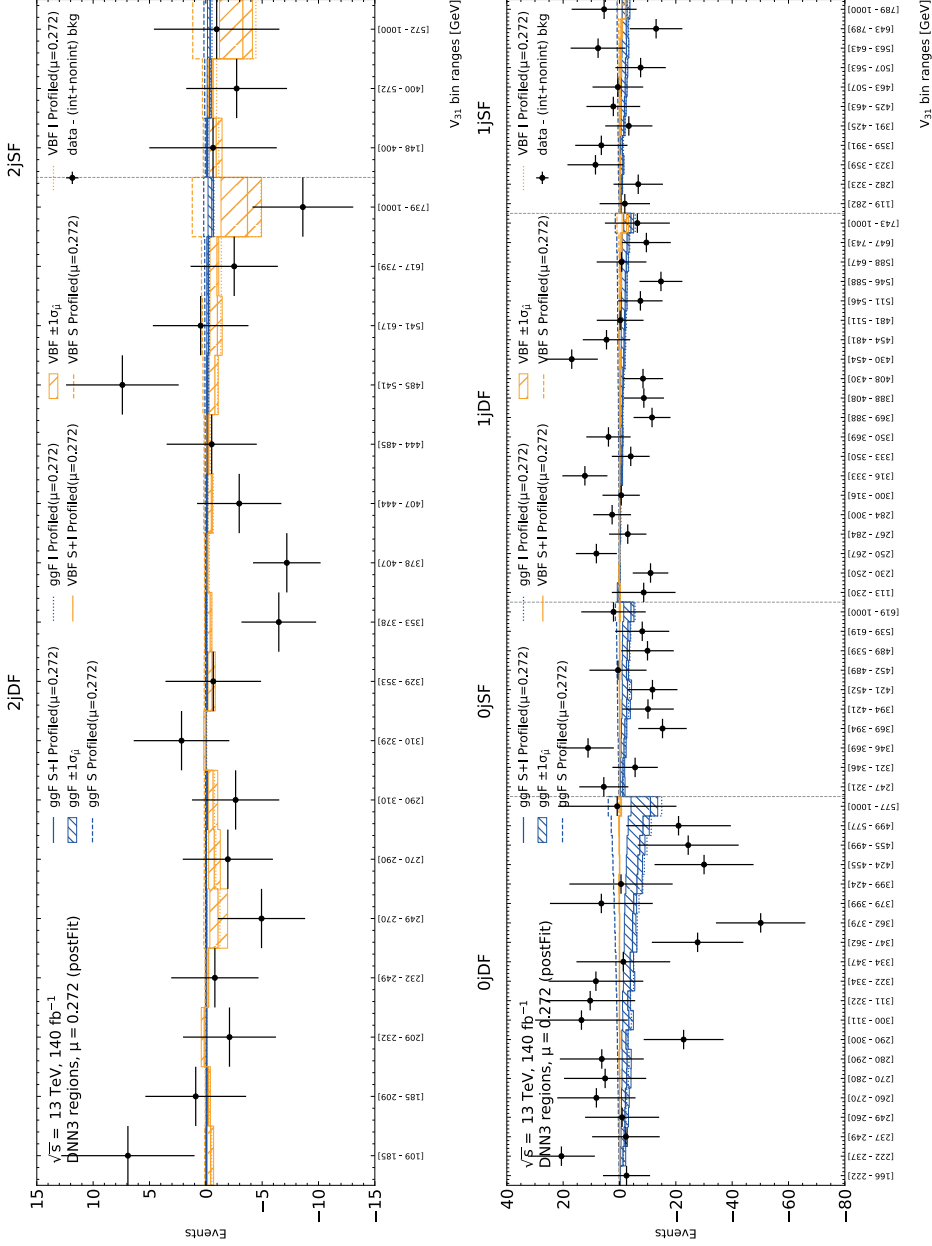


**Figure 10.29:** The post-fit plots of the  $\mu_{\text{off-shell}}$  signal strength measurement in the signal-enriched DNN regions. The y-axis is the density of events in each bin and the x-axis is the  $V_{0.3|1.0}$  variable. The goodness of fit metrics are shown in the top of each plot region along with the significance of the data points in each of the bins in grey. The uncertainty in the top and ratio pad corresponds to the total post-fit theoretical, experimental and statistical uncertainty.





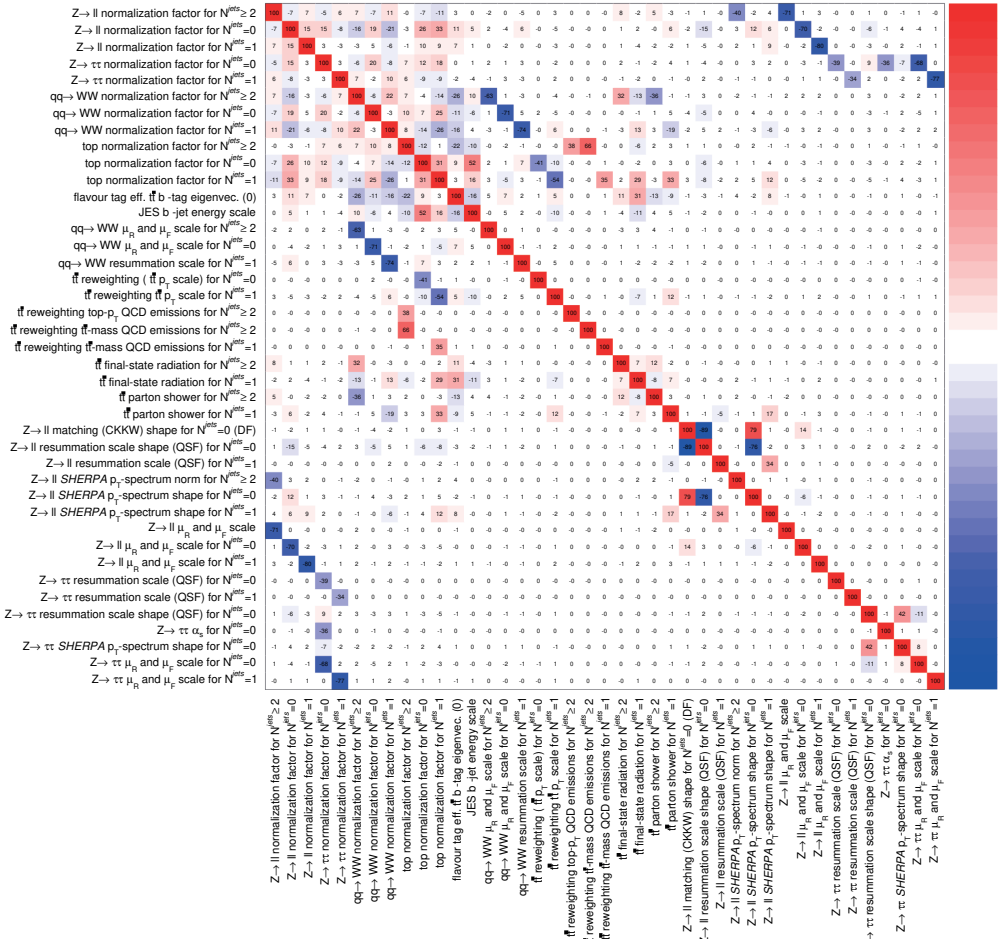
**Figure 10.31:** The post-fit signal residuals of the  $\mu_{\text{off-shell}}$  for the ggF and VBF signal strength measurement for each of the jet multiplicity regions in the mixed DNN regions. The residuals are shown in terms of the density of events in each bin and split into their interfering (S+I) and pure signal components. The uncertainty in the top and ratio pad corresponds to the total post-fit theoretical, experimental and statistical uncertainty.



**Figure 10.32:** The post-fit signal residuals of the  $\mu_{\text{off-shell}}$  for the ggF and VBF signal strength measurement for each of the jet multiplicity regions in the signal-enriched DNN regions. The residuals are shown in terms of the density of events in each bin and split into their interfering (S+I) and pure signal components. The uncertainty in the top and ratio pad corresponds to the total post-fit theoretical, experimental and statistical uncertainty.

### 10.10.1 Breakdown of uncertainties

The observed correlation between parameters  $X$  and  $Y$ ,  $\rho_{X,Y}$ , are calculated from the covariance matrix as  $\rho_{X,Y} = \text{cov}_{X,Y} / \sqrt{\text{cov}_{X,X}\text{cov}_{Y,Y}}$ , and are shown in Fig. 10.33.



**Figure 10.33:** The observed correlations exceeding 30% of the  $\mu_{\text{off-shell}}$  fit, which involve experimental uncertainties, theoretical uncertainties and background normalizations.

In order to understand the impact of the uncertainties on the off-shell measurement, the effect of each systematic uncertainty source on the parameters of interest is evaluated. The impact of a particular systematic uncertainty is quantified by the change of best-fit value of the POI when the uncertainty is varied by  $\pm 1\sigma$ . By minimizing the profile likelihood function with and without shifting the nuisance parameter value and comparing the best-fit values, the up and down impacts of the systematic uncertainty,  $\Delta\mu^{\text{up/down}}$ , can be

quantified as

$$\Delta\mu^{\text{up/down}} = \sqrt{\sigma_{\hat{\mu}}^2 - \sigma_{\hat{\mu}'}^2} (\text{NP} \pm 1\sigma),$$

where  $\sigma_{\hat{\mu}}^2$  and  $\sigma_{\hat{\mu}'}^2$  are the standard deviations of the best-fit values of the POI with and without the systematic uncertainty varied by  $\pm 1\sigma$ , respectively. The relative impact of the different sources of systematic uncertainty vary over the specific event topology. Tab. 10.12 shows the impact on the total nominal uncertainty of all systematic uncertainty categories on  $\mu_{\text{off-shell}}$ . The largest source of uncertainty comes from the data statistics, followed by the theory uncertainty and experimental uncertainties.

**Table 10.12:** Breakdown of observed impacts of uncertainty sources on the POI  $\mu_{\text{off-shell}}$  across all channels. The values represent the relative uncertainty  $\sqrt{\sigma_{\hat{\mu},\text{nom}}^2 - \sigma_{\hat{\mu},\text{NP}=\hat{\theta}}^2} / \sigma_{\hat{\mu},\text{nom}}$ , where  $\sigma_{\hat{\mu},\text{nom}}$  is the nominal uncertainty with all nuisance parameters floating, and  $\sigma_{\hat{\mu},\text{NP}=\hat{\theta}}$  is the uncertainty with specific parameters fixed to their best-fit values.

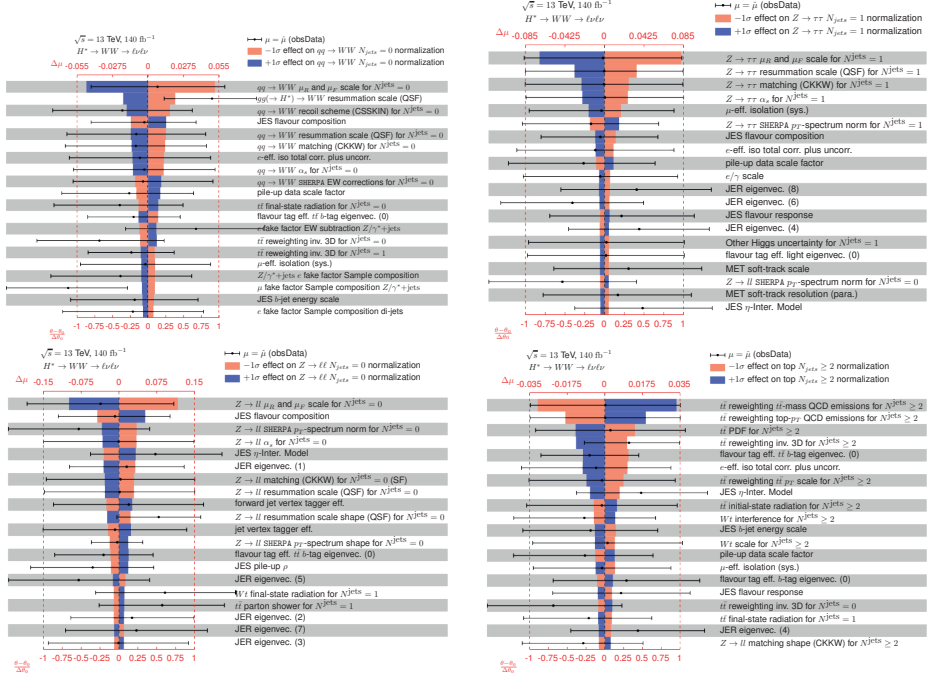
|                            |       |
|----------------------------|-------|
| Observed stat. uncertainty | 54.8% |
| MC stat. uncertainty       | 15.6% |
| Theory uncertainty         | 40.3% |
| Theory background          | 22.8% |
| Theory signal              | 34.8% |
| Experimental uncertainty   | 24.8% |
| Jets                       | 19.8% |
| Leptons                    | 6.31% |
| Others                     | 8.22% |
| Background normalisation   | 8.68% |
| MisID rate                 | 3.94% |

In order to understand which systematic uncertainty sources most limits the measurement, the top ranking uncertainties impacting the most important background normalizations are shown in Fig. 10.34; this illustrates the various sources limiting the background prediction. In addition to these, the top ranking nuisance parameters on  $\mu_{\text{off-shell}}$ ,  $\mu_{\text{off-shell}}^{\text{ggF}}$  and  $\mu_{\text{off-shell}}^{\text{VBF}}$  are shown in Figs. 10.35 to 10.37 to highlight the most important sources contributing to the uncertainty of the individual POIs.

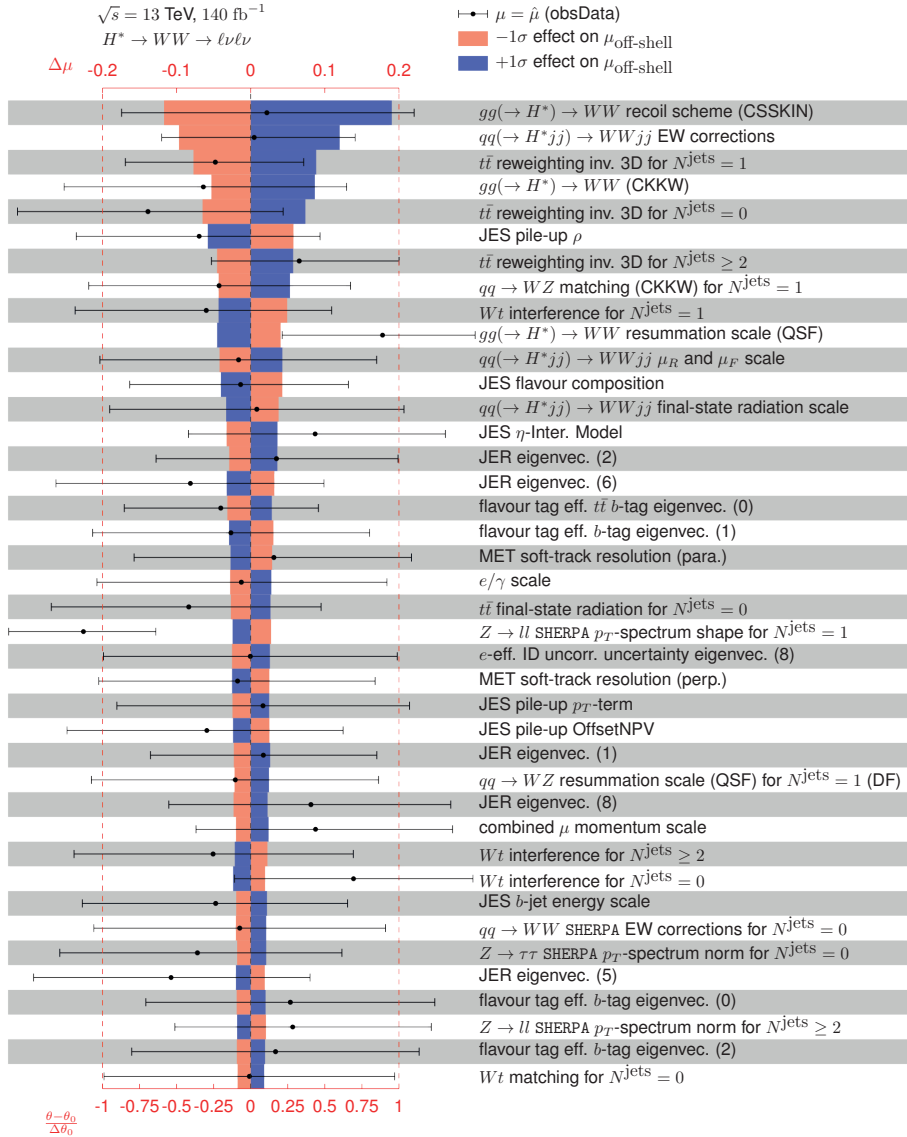
In these figures, all constrained NPs are intrinsically expressed in pull form, where nominally a measured value of 0 is expected. The NP is defined such that a pull of  $\pm 1$  corresponds to a variation of  $\pm 1\sigma$  of the nuisance parameter's uncertainty it represents, which makes it easy to compare the observed fit result with the expectation. The post-fit uncertainty of the NP is determined via MINOS and its correlations with the other NPs and the POIs are evaluated via HESSE and its diagonal elements of the covariance matrix. The largest sources of uncertainty are the  $gg \rightarrow WW$  theory and the VBF theory uncertainties, which are expected to be improved in the near future. These are followed by the  $t\bar{t}$  modelling uncertainty and the JES uncertainties. The most impactful uncertainty



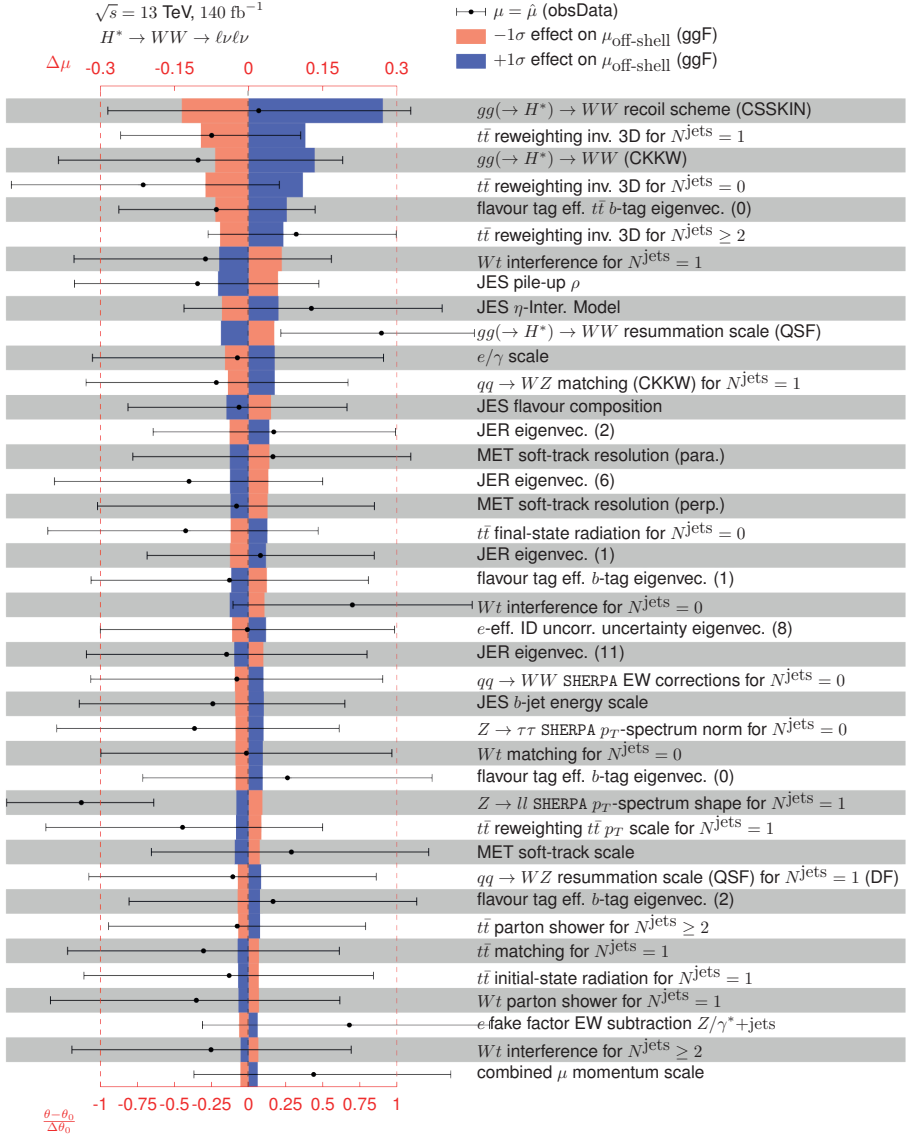
is the ggF recoil scheme shower uncertainty, which introduces large shape and migration uncertainties in the  $m_{WW}$  proxy and the signal categorisation. More discussion on the improvement of uncertainties can be found in [Part V](#).



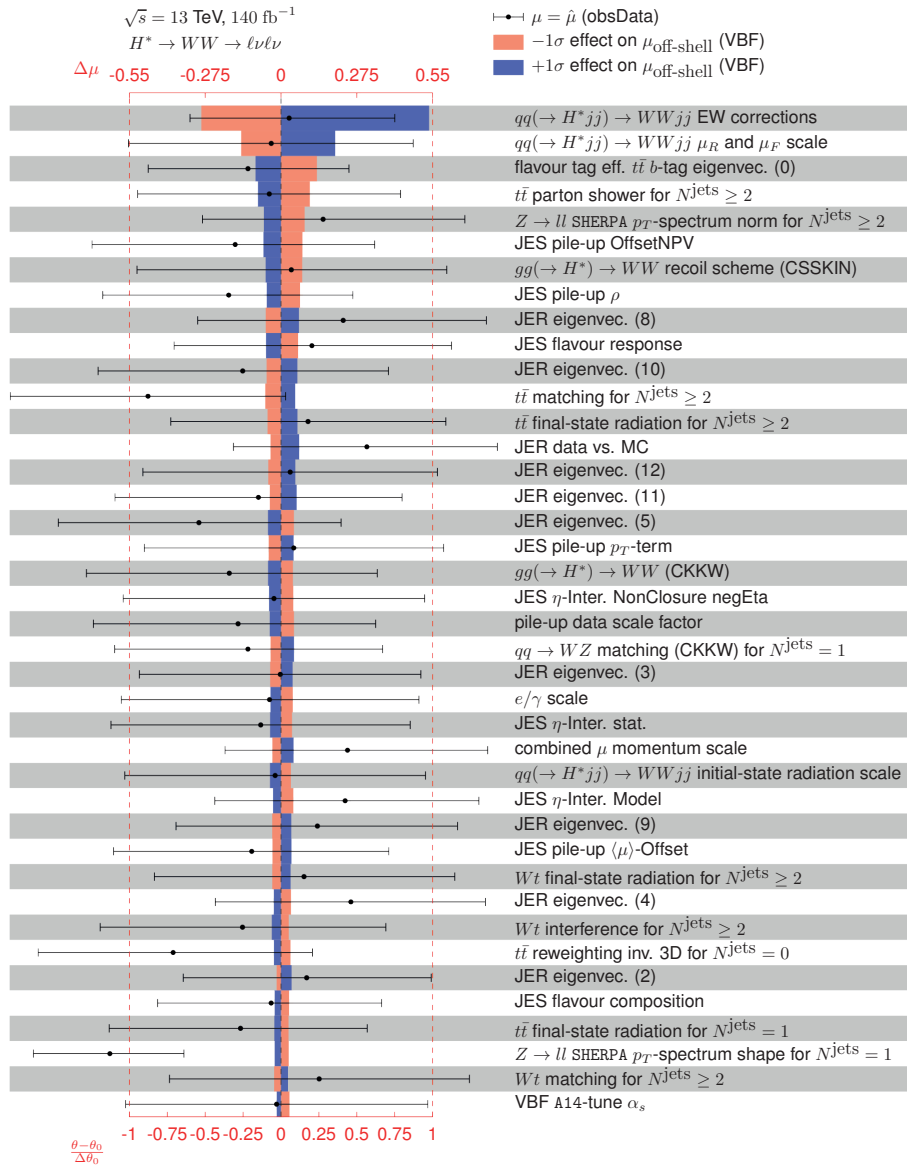
**Figure 10.34:** The impact of the top 20 systematic uncertainties on  $qq \rightarrow WW$ , Top,  $Z \rightarrow \tau\tau$  and  $Z \rightarrow \ell\ell$  normalization factors, sorted in decreasing order. The orange and purple bars represent the  $\pm 1\sigma$  shifts of the systematic uncertainties. The black data points and error bars represent the best-fit value and the total uncertainty of the nuisance parameter.



**Figure 10.35:** The impact of the systematic uncertainties on the  $\mu_{\text{off-shell}}$  signal strength measurement, sorted in decreasing order. The orange and purple bars represent the  $\pm 1\sigma$  shifts of the systematic uncertainties. The black data points and error bars represent the best-fit value and the total uncertainty of the nuisance parameter.



**Figure 10.36:** The impact of the systematic uncertainties on the  $\mu_{\text{off-shell}}$  ggF signal strength measurement, sorted in decreasing order. The orange and purple bars represent the  $\pm 1\sigma$  shifts of the systematic uncertainties. The black data points and error bars represent the best-fit value and the total uncertainty of the nuisance parameter.

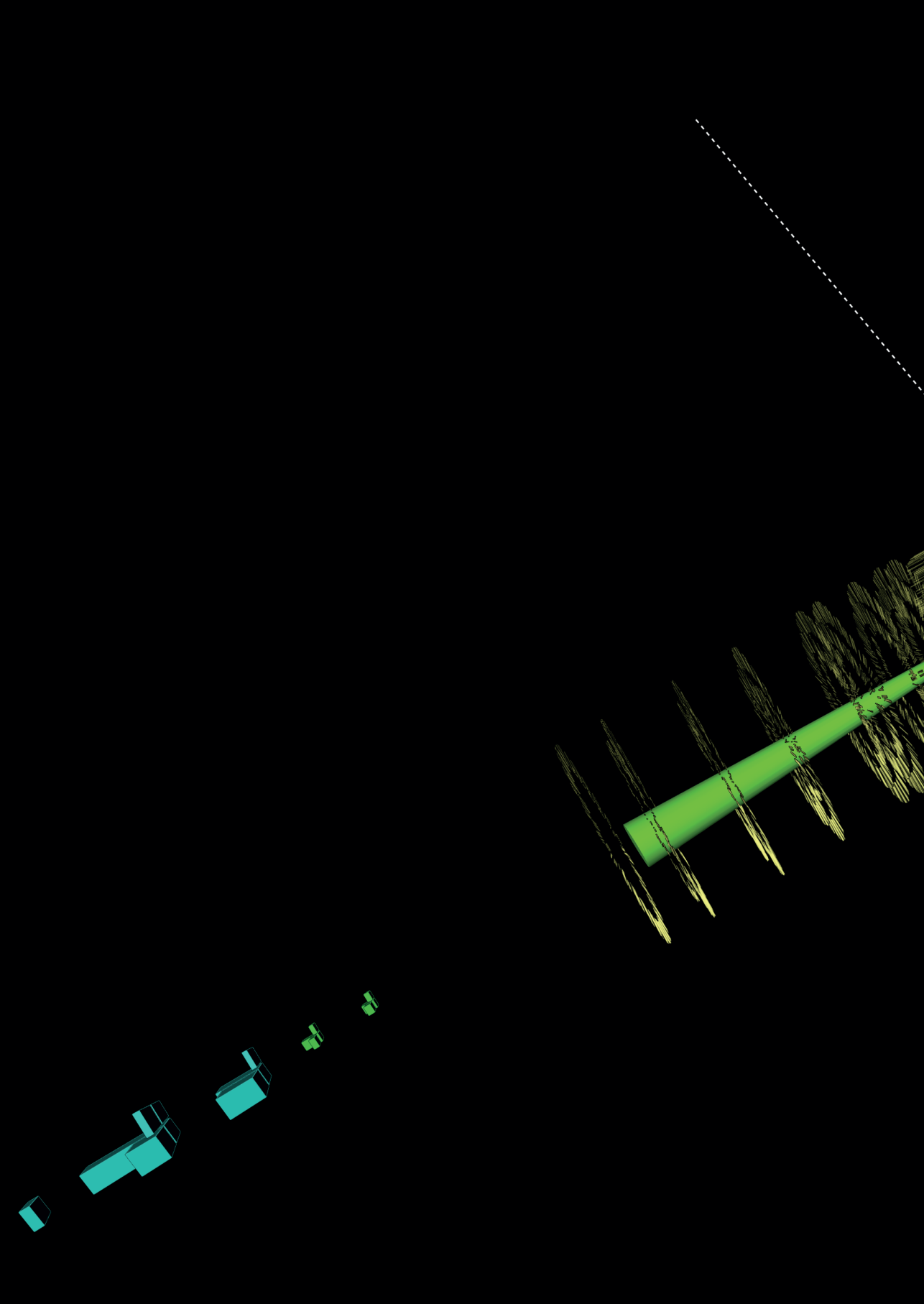


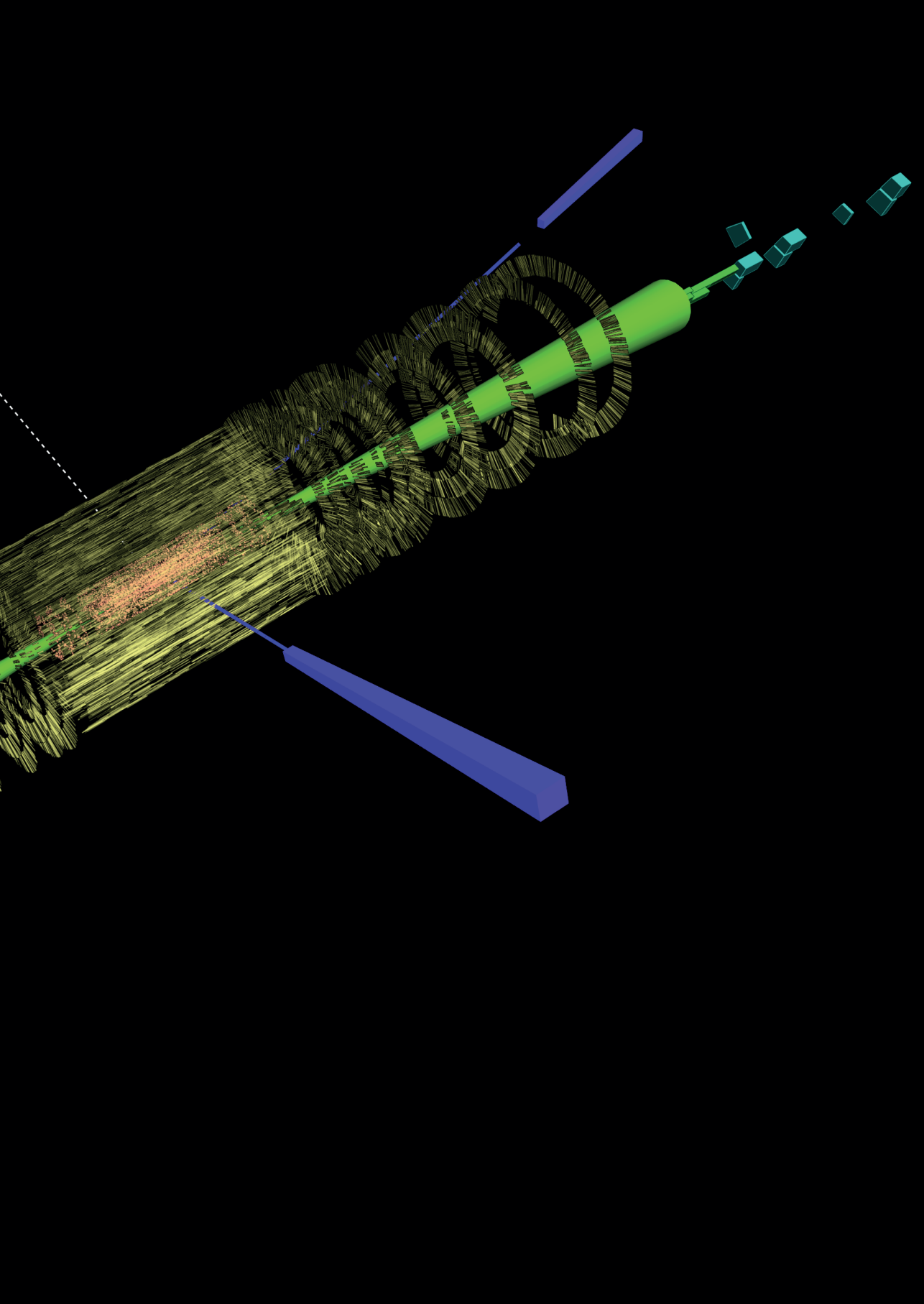
**Figure 10.37:** The impact of the systematic uncertainties on the  $\mu_{\text{off-shell}}$  VBF signal strength measurement, sorted in decreasing order. The orange and purple bars represent the  $\pm 1\sigma$  shifts of the systematic uncertainties. The black data points and error bars represent the best-fit value and the total uncertainty of the nuisance parameter.



PART IV

COMBINATIONS









## COMBINED ON-SHELL AND OFF-SHELL $H \rightarrow WW$ MEASUREMENT OF THE HIGGS TOTAL WIDTH

In this chapter, the on-shell and off-shell  $H \rightarrow WW$  analyses are combined to measure the Higgs boson width [4, 5]. Only the necessary details of the on-shell analysis, needed to perform the combination, are described in this chapter.

### 11.1 Data and input measurements

The same dataset used for the off-shell analysis is used for its on-shell counterpart, the full Run-2 dataset of  $140 \text{ fb}^{-1}$ . The triggers used are detailed in Tab. 6.2 and are identical with respect to the off-shell analysis, except for the addition of a di-lepton trigger used to populate the lower- $p_T$  spectra of the leading-leptons, which is useful for the on-shell low- $m_{\ell\ell}$  signature.

### 11.2 The on-shell $H \rightarrow WW^*$ measurement

The specific on-shell  $H \rightarrow WW^*$  analysis used for the Higgs boson lifetime measurement is the Run-2 legacy  $H \rightarrow WW^*$  measurement [5]. This measurement targets the on-shell ggF and VBF Higgs production mode in the  $H \rightarrow WW^*$  decay channel via the use of the  $e\mu/\mu e$   $N_{\text{jet}} = 0$ ,  $N_{\text{jet}} = 1$  and  $N_{\text{jet}} \geq 2$  channels and the  $N_{\text{jet}} \geq 2$   $ee/\mu\mu$  channel via the selections in Tab. 11.1.

**Table 11.1:** Event selection criteria used to define the SRs in the on-shell analysis.

| $N_{\text{jet}} = 0 \text{ } e\mu/\mu e$                                                                                                                                                                                                                                                                                                                                                   | $N_{\text{jet}} = 1 \text{ } e\mu/\mu e$ | $N_{\text{jet}} \geq 2 \text{ } e\mu/\mu e/ee/\mu\mu \text{ ggF-like}$                                                    | $N_{\text{jet}} \geq 2 \text{ } e\mu/\mu e/ee/\mu\mu \text{ VBF-like}$                       |
|--------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------------|------------------------------------------|---------------------------------------------------------------------------------------------------------------------------|----------------------------------------------------------------------------------------------|
| Preselection                                                                                                                                                                                                                                                                                                                                                                               |                                          |                                                                                                                           |                                                                                              |
| Two isolated, different-flavour (same-flavour, for $ee/\mu\mu$ channels) leptons with opposite charge<br>$p_{T,\text{lead}} > 22 \text{ GeV}$ and $p_{T,\text{sublead}} > 15 \text{ GeV}$<br>$m_{\ell\ell} > 10 \text{ GeV}$ (12 GeV for $ee/\mu\mu$ channels)<br>$(\mathcal{S}_{\text{miss}} > 4 \text{ and } p_{\text{T}}^{\ell\ell} > 40 \text{ GeV, for } ee/\mu\mu \text{ channels})$ |                                          |                                                                                                                           |                                                                                              |
| Background rejection                                                                                                                                                                                                                                                                                                                                                                       |                                          |                                                                                                                           |                                                                                              |
| $n_{b\text{-jets}}(p_T > 20\text{GeV}) = 0$                                                                                                                                                                                                                                                                                                                                                |                                          |                                                                                                                           |                                                                                              |
| $\Delta\phi(\ell, \ell + \text{MET}) > 1.57 \text{ rad}$<br>$\Delta R(\ell, \ell) > 0.6$                                                                                                                                                                                                                                                                                                   | $m_{\text{T}} < m_Z - 25 \text{ GeV}$    |                                                                                                                           |                                                                                              |
| $H^* \rightarrow WW \rightarrow \ell\nu\ell\nu$ topology                                                                                                                                                                                                                                                                                                                                   |                                          |                                                                                                                           |                                                                                              |
| MET > 20 GeV<br>$\Delta\phi(\ell, \ell) < 2.0 \text{ rad}$<br>$m_{\ell\ell} < 55 \text{ GeV}$                                                                                                                                                                                                                                                                                              | -                                        | $\Delta\phi(\ell, \ell) < 1.8 \text{ rad}$<br>$m_{\ell\ell} < 55 \text{ GeV}$<br>VH orthogonality<br>Fail CJV or Fail OLV | -<br><br>$m_{\ell\ell} < 70 \text{ GeV}$<br>$m_{\text{jj}} > 120 \text{ GeV}$<br>CJV and OLV |

The measurements are designed to be mutually exclusive, such that no overlap is present between the off-shell and on-shell channels, allowing for a straightforward combination. Both of the analyses use the same background estimation techniques and only differ where the signal samples are concerned. This allows for the correlation of theory uncertainties between the two, allowing for the cancellation of systematic uncertainties in the combination.

In the combination, the two analysis share the on-shell  $N_{\text{jet}} = 1$   $Z \rightarrow \tau\tau$  control region with  $60 < p_T^H < 120$  GeV, which is similar in phase-space and obtained normalization factor when compared to the  $N_{\text{jet}} = 1$   $Z \rightarrow \tau\tau$  control region used in the off-shell stand-alone measurement.

Additionally, the on-shell analysis uses a dedicated  $WW$   $N_{\text{jet}} = 1$  control region designed to be orthogonal to the off-shell measurement. The orthogonal  $WW$   $N_{\text{jet}} = 1$  control region is defined via a  $55 < m_{\ell\ell} < 80$  GeV, a  $\Delta\phi(\ell, E_T^{\text{miss}})/\Delta\phi_{\ell\ell} > 1.25$  and a  $m_T > 110$  GeV requirement, which increases the  $WW$  purity of the control region by filtering out  $Z \rightarrow \tau\tau$  events.

The Run-2 legacy  $H \rightarrow WW^*$  analysis [5] provides the following on-shell signal strength measurements for the ggF and VBF production modes:

$$\begin{aligned}\mu_{\text{on-shell}}^{\text{ggF}} &= 1.18_{-0.07}^{+0.07} (\text{stat.})_{-0.11}^{+0.13} (\text{syst.})_{-0.13}^{+0.14} (\text{tot.}) \\ \mu_{\text{on-shell}}^{\text{VBF}} &= 0.95_{-0.13}^{+0.13} (\text{stat.})_{-0.14}^{+0.18} (\text{syst.})_{-0.19}^{+0.22} (\text{tot.}),\end{aligned}$$

### 11.2.1 Correlation of systematic uncertainties

As the two analyses are performed in the same analysis framework and use the same background estimation techniques, all experimental uncertainty sources used in the analyses are identical and are therefore correlated. Secondly, the same misidentified lepton background estimation is used; its uncertainties are therefore treated as correlated. Additionally, the theoretical uncertainty sources for the background are correlated between the two analyses where possible, with only a few exceptions.

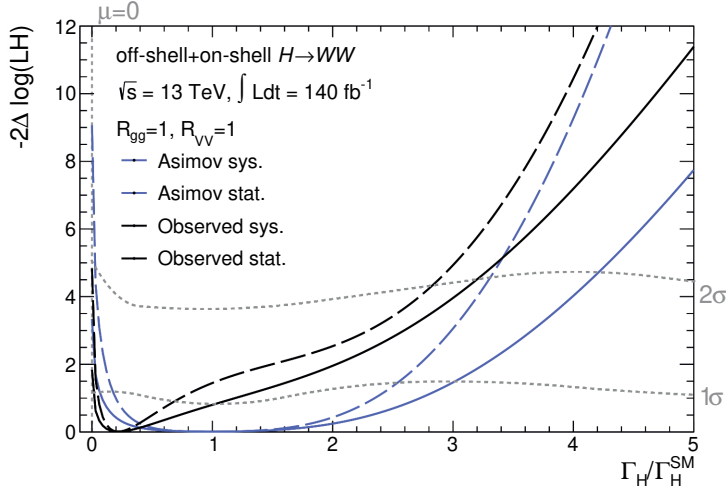
As the VBF on-shell analysis treats EW  $qq \rightarrow WWjj$  as a background, these theory uncertainties are correlated with the signal uncertainties of the off-shell analysis. The remaining theory uncertainties on the signal are left uncorrelated, as they cover very different phase-spaces and uncertainty sources. This approach is deemed conservative as in this configuration the uncertainties do not cancel out in the combination.

## 11.3 Results

The results in this section are split into the width and off-shell/on-shell coupling ratio measurements. Both statistical models are built using different assumptions and provide different off-shell Higgs boson property information.

### 11.3.1 Higgs width and lifetime results

Using the same likelihood model as the off-shell analysis in combination with the on-shell analysis, a likelihood scan of the Higgs boson width, shown in Fig. 11.1, is performed via the parameterization explained in § 3.4.3.



**Figure 11.1:** The likelihood scan of the Higgs boson width measured in the  $H^* \rightarrow WW$  channel by combining the on-shell and off-shell analyses, along with the  $1\sigma$  and  $2\sigma$  constructed confidence intervals.

The likelihood scan results in a best-fit value of

$$\Gamma_H/\Gamma_H^{\text{SM}} = 0.22^{+0.47}_{-0.17} (\text{stat.})^{+0.67}_{-0.11} (\text{syst.})^{+0.82}_{-0.21} (\text{tot.}).$$

The measurement of  $\Gamma_H/\Gamma_H^{\text{SM}}$  is translated into the Higgs boson width using the SM width of  $\Gamma_H^{\text{SM}} = 4.1$  MeV and results in a measured width  $\Gamma_H$  and corresponding lifetime  $\tau_H = \hbar/\Gamma_H$  of

$$\begin{aligned} \Gamma_H &= 0.90^{+1.93}_{-0.72} (\text{stat.})^{+2.75}_{-0.45} (\text{syst.})^{+3.36}_{-0.85} (\text{tot.}) \text{ MeV} \\ \tau_H &= 7.3^{+86.7}_{-5.7} \times 10^{-22} \text{ s (tot.)}. \end{aligned}$$

The observed (expected) Higgs width is in agreement with the SM prediction, with a  $p$ -value of 0.23 (0.11) estimated from a distribution of the test statistic  $q_0$  obtained from toy ensemble studies. With respect to the background-only hypothesis, the observed (expected)  $\Gamma_H$  corresponds to an observation (expectation) significance  $Z$  of  $0.7\sigma$  ( $1.2\sigma$ ). An observed (expected) upper limit on the total Higgs width of  $\Gamma_H$  of 13.1 (17.3) MeV is found, at 95% CL<sup>1</sup>.

<sup>1</sup> This 95% CL upper limit is derived using the  $t_\mu$  test statistic, which considers fluctuations in both directions. This will overcover when interpreted as a one-sided limit. The quoted limited limit is therefore somewhat conservative.

The confidence intervals and uncertainties for the results of the combined on-shell and off-shell analysis are calculated by following the procedure of Neyman construction as explained in § 10.9.2. Fig. 11.2 shows the full confidence belt for  $\mu = [0, 8]$ , which represents the signal strength scaling the observed Higgs width  $\mu \equiv \Gamma_H / \Gamma_H^{SM}$ .

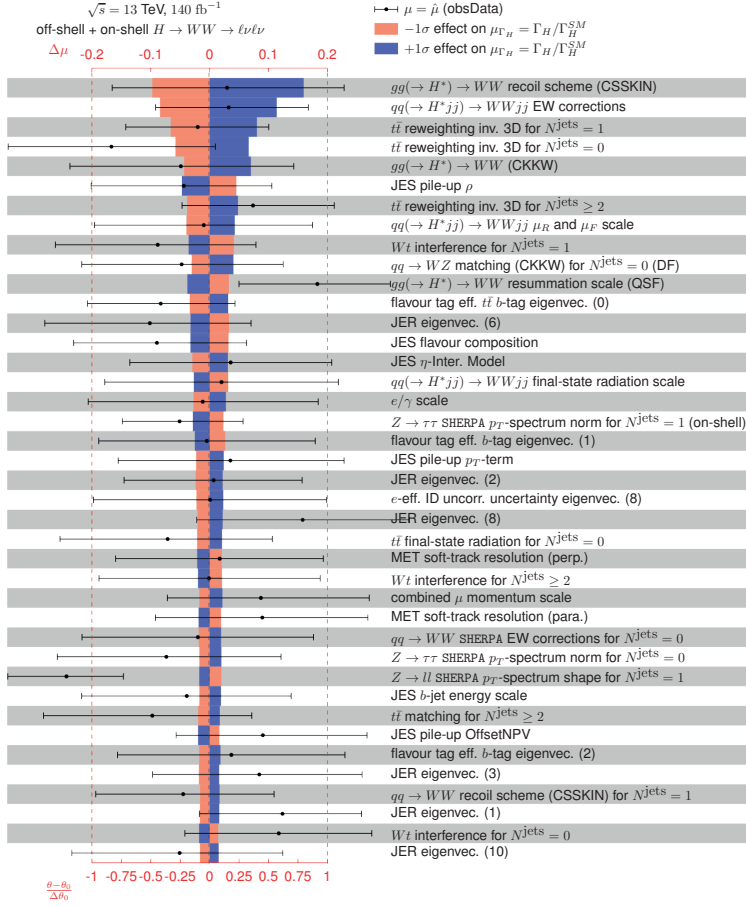


**Figure 11.2:** The critical values of the test statistic for the POI derived from the Neyman construction procedure compared to the asymptotic approximation. The black lines represent the estimated critical values along with their uncertainty, the yellow and green band borders are determined via a polynomial fit to the data as function of the injected, true, value of  $\mu$ , with which the toys were generated. The parameter  $\mu$  now represents the Higgs width normalized to the SM,  $\Gamma_H / \Gamma_H^{SM}$ .

An impact summary plot of the systematic uncertainties, broken down in terms of source on the width measurement, is shown in Fig. 11.3. It is observed that the top ranking systematic uncertainties are dominated by the same sources as the off-shell analysis, as well as the fact that the lower limit is statistically limited and the upper limit is systematically limited. This is expected, since the off-shell measurement sensitivity is limiting the width measurement most significantly and the on-shell VBF and ggF measurements are systematically limited.

### 11.3.2 Off-shell/on-shell coupling ratio results

When a SM Higgs width is assumed, a measurement of the off-shell/on-shell coupling ratio can be performed. By simultaneously profiling the vector boson coupling ratio and the gluon fusion coupling ratio, a 2D likelihood scan of the off-shell/on-shell coupling ratios is determined, shown in Fig. 11.4. The 2-POI best-fit values are  $R_{gg} = 0.22$  and  $R_{VV} = 0.63^{+1.57}_{-0.59}$ .

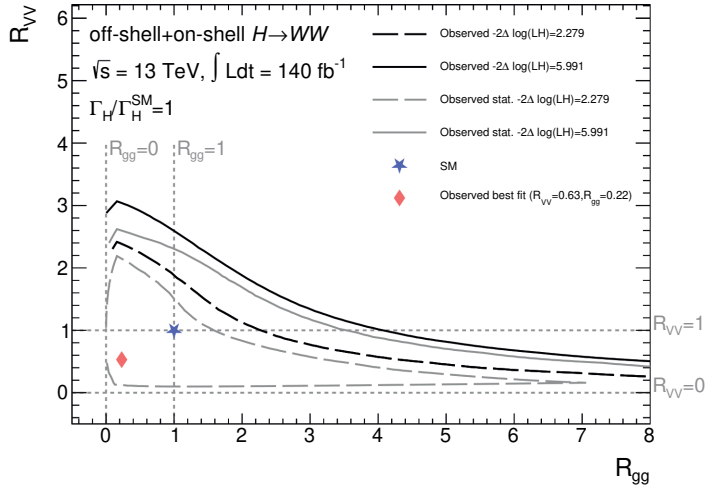


**Figure 11.3:** The impact of the systematic uncertainties on the Higgs boson width measurement sorted in decreasing order. The orange and purple bars represent the  $\pm 1\sigma$  shifts of the systematic uncertainties. The black data points and error bars represent the best-fit value and the total uncertainty of the nuisance parameter.

From the 2-POI fit, an observed (expected) limit of  $R_{VV} < 2.8$  (2.7) at 95% CL is found<sup>2</sup>. The confidence intervals of the  $R_{VV}$  measurement and its likelihood scan, shown in Fig. 11.5, are calculated by following the procedure of Neyman construction. Fig. 11.6 shows the full confidence belt for  $\mu = [0, 4]$ , which represents the off-shell/on-shell  $VV$  coupling ratio  $\mu \equiv R_{VV}$ .

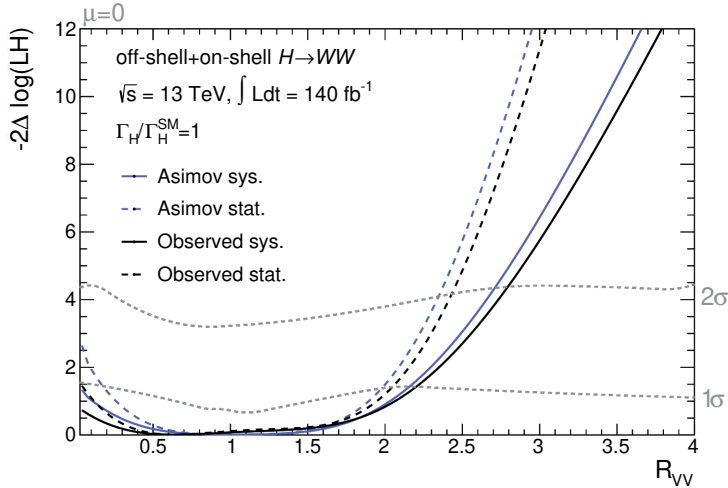
By decoupling the production and decay into  $R_{gg}$  and  $R_{VV}$ , a flat direction in the likelihood is added and sensitivity is lost. Because of this new degeneracy in the fit, the analysis is unable to set well-defined 1- and 2 $\sigma$  limits on the parameters.

<sup>2</sup>This 95% CL upper limit is derived using the  $t_\mu$  test statistic, which considers fluctuations in both directions. This will overcover when interpreted as a one-sided limit. The quoted limited limit is therefore somewhat conservative.

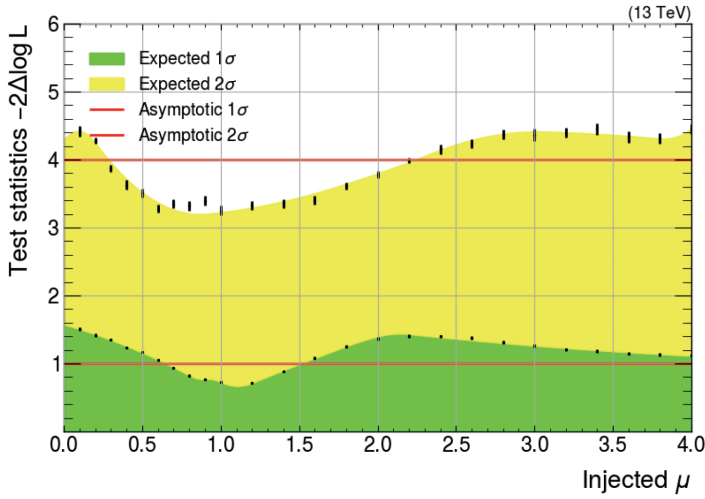


**Figure 11.4:** The 2D likelihood scan of the observed off-shell/on-shell coupling ratios in the  $H^* \rightarrow WW$  channel showing the  $1\sigma$  and  $2\sigma$  contours of the asymptotic confidence intervals. This scan is performed under the assumption of a SM Higgs boson width.

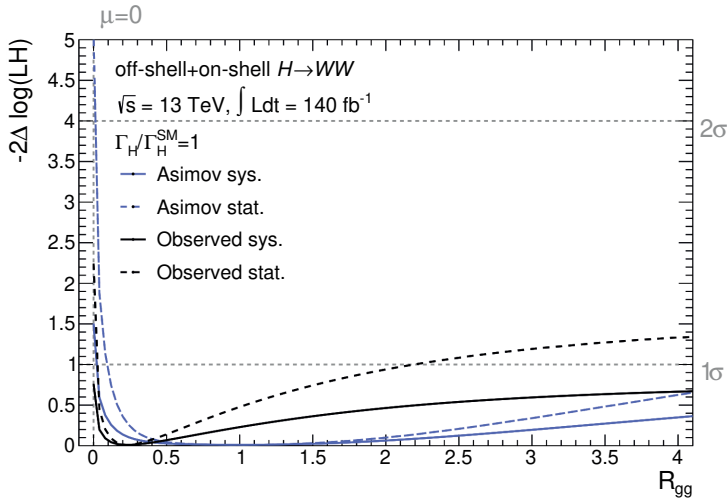
The flat behaviour of the  $R_{gg}$  likelihood scan in Fig. 11.7 is explained by the fact that  $R_{VV}$  can directly compensate for any increase in  $R_{gg}$ , which keeps the total yield low. This indicates little sensitivity towards low values of  $R_{VV}$  in combination with high values of  $R_{gg}$ .



**Figure 11.5:** The 1D likelihood scan of the off-shell/on-shell vector boson coupling ratio. This scan is performed under the assumption of a SM Higgs boson width and by simultaneously profiling the gluon fusion coupling ratio. The constructed  $1\sigma$  and  $2\sigma$  confidence intervals are shown.



**Figure 11.6:** The critical values of the test statistic for the  $R_{VV}$  POI derived from the Neyman construction procedure compared to the asymptotic approximation. The black lines represent the estimated critical values along with their uncertainty, the yellow and green band borders are determined via a polynomial fit to the data as function of the injected, true, value of  $\mu$ , with which the toys were generated.



**Figure 11.7:** The 1D likelihood scan of the off-shell/on-shell gluon coupling ratio. This scan is performed under the assumption of a SM Higgs boson width and by simultaneously profiling the vector boson coupling ratio. The asymptotic 1 $\sigma$  and 2 $\sigma$  confidence intervals are shown.





## COMBINED SM-EFT INTERPRETATION OF HIGGS BOSON PRODUCTION AND DECAY

This chapter presents an interpretation of the results from the combined Higgs boson measurements in production and decay using differential STXS measurements [6], within the framework of the SM effective field theory. The global signal strength measurement of this analysis combines the  $H \rightarrow \gamma\gamma, ZZ^*, WW^*, \tau\tau, b\bar{b}, \mu\mu, Z\gamma$  processes and searches for decays into invisible final states using up to  $139 \text{ fb}^{-1}$  of  $pp$ -collision data collected at  $\sqrt{s} = 13 \text{ TeV}$  with the ATLAS detector. The STXS measurement of the same analysis combines the channels  $H \rightarrow \gamma\gamma, ZZ^*, WW^*, \tau\tau, b\bar{b}$ .

### 12.1 Input measurements for the combination

The analysed decay channels, targeted production modes and integrated luminosities of the datasets from each input analysis entering the combined measurement are shown in Tab. 12.1.

**Table 12.1:** The decay channels, targeted production modes and integrated luminosity ( $\mathcal{L}$ ) used for each input analysis of the combination. The analysis channels are used for both the STXS and SM-EFT interpretations.

| Decay channel                | Target Production Modes                                                 | $\mathcal{L} [\text{fb}^{-1}]$ | Ref.      |
|------------------------------|-------------------------------------------------------------------------|--------------------------------|-----------|
| $H \rightarrow \gamma\gamma$ | ggF, VBF, $WH$ ,<br>$ZH, t\bar{t}H, tH$                                 | 139                            | [240]     |
| $H \rightarrow ZZ^*$         | ggF, VBF, $WH$<br>$ZH, t\bar{t}H(4\ell)$                                | 139                            | [241]     |
| $H \rightarrow WW^*$         | ggF, VBF                                                                | 139                            | [242]     |
| $H \rightarrow \tau\tau$     | ggF, VBF, $WH$ ,<br>$ZH, t\bar{t}H(\tau_{\text{had}}\tau_{\text{had}})$ | 139                            | [243]     |
| $H \rightarrow b\bar{b}$     | $WH, ZH$                                                                | 139                            | [244–246] |
|                              | VBF                                                                     | 126                            | [247]     |
|                              | $t\bar{t}H$                                                             | 139                            | [248]     |

Further details about the analysis procedure and reconstruction of the final states can be found in the analysis references in the table. Before the analyses are combined, it has

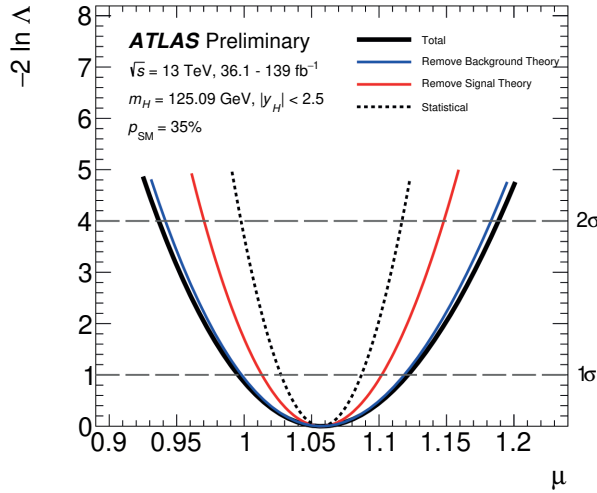
been checked if these analyses are using a statistically independent dataset, which is needed in order to avoid double counting of events.

## 12.2 Inclusive and STXS style results

Although this combination is performed using the STXS framework, it can be simplified to the inclusive signal strength measurement. By combining all of STXS categories, the global Higgs signal strength is measured to be

$$\begin{aligned}\mu_{|y_H| < 2.5}^{pp \rightarrow H} &= 1.06 \pm 0.03 \text{ (stat.)} \pm 0.03 \text{ (exp.)} \pm 0.04 \text{ (sig. theo.)} \pm 0.04 \text{ (bkg. theo.)} \\ &= 1.06 \pm 0.06 \text{ (tot.)}.\end{aligned}$$

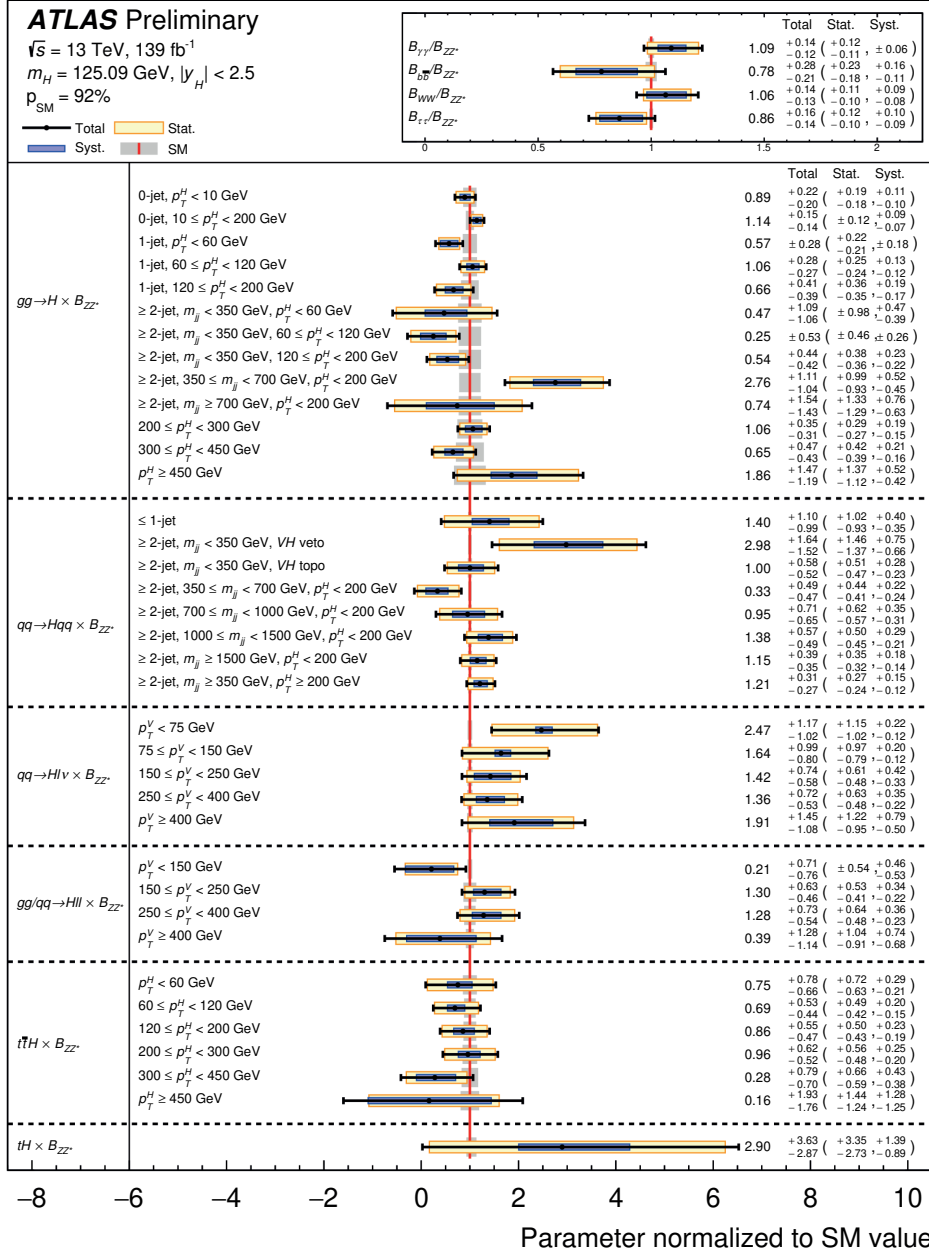
The total uncertainty is decomposed into statistical, experimental and separate theoretical uncertainties on the signal and background, with the likelihood scan of the observed data shown Fig. 12.1. Compared to the previous combination [249], the total measurement un-



**Figure 12.1:** Scan of the global Higgs signal strength with all systematic uncertainties profiled (solid black line). It also displays results without considering theoretical uncertainties in background processes (solid blue line) and those in the Higgs process (solid red line), as well as results based only on statistical uncertainties (dotted black line). The dashed horizontal lines represent the  $1\sigma$  and  $2\sigma$  confidence intervals for  $\mu$ . The agreement between the observed global signal strength and the SM prediction corresponds to a p-value of 35% [6].

certainty is reduced by 10%. The reduction in total uncertainty is driven by the reduction in statistical uncertainty due to the increased amount of data and by including additional channels. Currently, the total uncertainty is dominated by theoretical uncertainty.

The SM-EFT interpretation performed in this chapter uses a modified version of the STXS stage-1.2 framework as defined in § 3.3.3, where some original bins have been merged due to low statistics. The results summarized in Fig. 12.2 are obtained with a profile likelihood to the data, where the cross-sections of all the STXS categories are determined



**Figure 12.2:** The best-fit values and uncertainties for cross-sections in each STXS region and the ratios of branching fractions  $B_f/B_{ZZ^*}$  for Higgs decay channels  $\gamma\gamma, WW^*, \tau\tau, b\bar{b}$ , normalized to SM predictions. Black error bars, blue boxes, and yellow boxes indicate total, systematic, and statistical uncertainties, respectively. Grey bands represent theoretical uncertainties. The combined measurements agree with the SM prediction, corresponding to  $p_{\text{SM}} = 92\%$ .

simultaneously, normalized to their SM predictions. Their branching ratios, normalized to the branching ratio of Higgs decaying into two neutral vector bosons  $B_{ZZ}$ , are also floating.

### 12.3 The linearized SM-EFT model

The SM-EFT model, as introduced in Chapter 4, can describe effects of possible BSM physics as a function of Wilson coefficients generated by new physics at a mass scale  $\Lambda$ . The theory is able to provide predictions for experimental observables in terms of expansions around  $c_i/\Lambda$ . This is achieved using a series of operators  $\mathcal{O}_i^{(D)}$ , which consist of gauge-invariant combinations of SM fields with a mass dimension greater than 4. The SM-EFT lagrangian

$$\mathcal{L}_{\text{SM-EFT}} = \mathcal{L}_{\text{SM}} + \sum_i \frac{c_i^{(5)}}{\Lambda} \mathcal{O}_i^{(5)} + \sum_i \frac{c_i^{(6)}}{\Lambda^2} \mathcal{O}_i^{(6)} + \sum_i \frac{c_i^{(7)}}{\Lambda^3} \mathcal{O}_i^{(7)} + \sum_i \frac{c_i^{(8)}}{\Lambda^4} \mathcal{O}_i^{(8)} + \dots,$$

reduces to the SM lagrangian  $\mathcal{L}_{\text{SM}}$  when  $\Lambda \rightarrow \infty$ . For the interpretation of Higgs data, this general formulation of the SM-EFT lagrangian is simplified by ignoring odd-dimensional operators, which introduce lepton and baryon number violation, and by restricting the expansion to its leading terms. This simplification yields the linearized SM-EFT model used in this thesis, consisting of the SM lagrangian augmented by the complete set of dimension-6 operators:

$$\mathcal{L}_{\text{SM-EFT}} = \mathcal{L}_{\text{SM}}^{(4)} + \sum_i \frac{c_i^{(6)}}{\Lambda^2} \mathcal{O}_i^{(6)}.$$

In this equation,  $\mathcal{L}_{\text{SM}}^{(4)}$  represents the usual renormalizable part of the SM lagrangian. The value of  $\Lambda$  is set to 1 TeV motivated by convention, theory consideration and that it is sufficiently far away from the electroweak scale  $v = \sqrt{2G_F}$  of around 246 GeV. Results for alternative values of the new-physics scale,  $\Lambda'$ , can be obtained by multiplying the dimension-6 Wilson coefficients by  $(\Lambda'/\Lambda)^2$ . Wilson coefficients should not exceed 1, within the range of validity of the EFT [115].

Scattering cross-sections calculated using  $\mathcal{L}_{\text{SM-EFT}}$  are expressed as the SM scattering amplitude  $\mathcal{M}_{\text{SM}}$  with the addition of modifiers induced by the dimension-6 operators. These operators can introduce new vertices, redefine the SM fields or shift the input parameters  $m_W$ ,  $m_Z$  and  $G_F$ . The total scattering cross-section is then expanded in terms of the Wilson coefficients  $c_i$ , similar to  $\mathcal{L}_{\text{SM-EFT}}$ , as

$$\begin{aligned} \sigma &\propto |\mathcal{M}_{\text{SM-EFT}}|^2 = |\mathcal{M}_{\text{SM}} + \mathcal{M}_{\text{EFT}}|^2 \\ &= |\mathcal{M}_{\text{SM}} + \sum_i \frac{c_i^{(6)}}{\Lambda^2} \mathcal{M}_i^{(6)}|^2 + \mathcal{O}(\Lambda^{-4}) \\ &= |\mathcal{M}_{\text{SM}}|^2 + \sum_i \frac{c_i^{(6)}}{\Lambda^2} 2\text{Re}(\mathcal{M}_i^{(6)} \mathcal{M}_{\text{SM}}^*) + \mathcal{O}(\Lambda^{-4}). \end{aligned}$$

In this expansion,  $\mathcal{M}_{\text{SM}}$  is the matrix element as it occurs in the SM, and  $\mathcal{M}_i^{(6)}$  are amplitude corrections containing the operators  $\mathcal{O}_i^{(6)}$ . The leading SM-EFT contribution is the interference of dimension-6 new-physics effects with the SM amplitudes. Since only dimension-6 terms are available in models at present, the set of effects occurring at  $O(\Lambda^{-4})$  cannot be simulated fully. For example, the effect  $\mathcal{M}_i^{(8)}\mathcal{M}_{\text{SM}}$  occurring at  $O(\Lambda^{-4})$  would be missing. Therefore, in order to obtain a complete basis in which the measurements can be interpreted, the expansion is truncated at  $O(\Lambda^{-2})$  and the  $O(\Lambda^{-4})$  are neglected.

Using this equation, the cross-sections of a given process  $p$  acquires a linear dependence on the Wilson coefficients and is written as  $\sigma_{\text{SM-EFT}}^p = \sigma_{\text{SM}}^p + \sigma_{\text{interference}}^p = \sigma_{\text{SM}}^p (1 + \sum_i A_i^p c_i)$ , where  $\sigma_{\text{SM}}^p$  is the SM cross-section. The linear effect of the operator associated with the Wilson coefficient  $c_i$  on the cross-section is  $A_i = \sigma_{\text{interference}}^p / \sigma_{\text{SM}}^p$ , relative to the SM cross-section.

As the SM-EFT contributions are only available at leading order in QED and next-to-leading-order in QCD, the parameterization of the total cross-section is computed as a relative correction to the SM prediction at the highest available perturbative order. This methodology assumes the effect of the SM-EFT operators on the cross-section relative to the SM prediction remains constant at each perturbative order. Since the Higgs boson production cross-section calculations are performed at NNLO in QCD. Therefore, the total SM-EFT cross-section is computed as

$$\sigma_{\text{SM-EFT}} = \sigma_{\text{SM}}^{(\text{N})\text{NLO}} \left( 1 + \frac{\sigma_{\text{interference}}^{(\text{N})\text{LO}}}{\sigma_{\text{SM}}^{(\text{N})\text{LO}}} \right).$$

The resonant Higgs boson process, which is the signal of this analysis, is simulated using the narrow width approximation. This allows the operator impact on the total cross-section to be factorized into the product of the production cross-section and branching ratio into the final state of interest. Following this factorization, the impact of the SM-EFT operators on the branching ratio  $\mathcal{B}$  is computed as

$$\mathcal{B}_{\text{SM-EFT}}^{H \rightarrow X} = \frac{\Gamma_{\text{SM-EFT}}^{H \rightarrow X}}{\Gamma_{\text{SM-EFT}}^H} = \mathcal{B}_{\text{SM}}^{H \rightarrow X} \left( \frac{1 + \sum_i c_i \Gamma_i^{H \rightarrow X}}{1 + \sum_i c_i \Gamma_i^H} \right) + O(\Lambda^{-4}).$$

In this equation,  $\sum_i c_i \Gamma_i^{H \rightarrow X}$  is the summed linear correction to the partial width of the final state  $X$ ,  $\Gamma^{H \rightarrow X}$  and  $\sum_i c_i \Gamma_i^H$  is the summed correction to the total Higgs boson width  $\Gamma^H$ . All the sums run over the Wilson coefficients  $i$  which affect the decay process. The linearized expression for the product of production and decay is given by:

$$\begin{aligned} \sigma_{\text{SM-EFT}}^p \cdot \mathcal{B}_{\text{SM-EFT}}^{H \rightarrow X} &= \sigma_{\text{SM}}^p \mathcal{B}_{\text{SM}}^{H \rightarrow X} \left( 1 + \sum_i A_i^p c_i + \sum_i c_i \Gamma_i^{H \rightarrow X} - \sum_i c_i \Gamma_i^H \right) \\ &+ O(\Lambda^{-4}). \end{aligned}$$

In the linearization, higher-order terms of the denominator are truncated, to obtain a description valid up to  $O(\Lambda^{-2})$ . The constant terms  $A_i^p$ ,  $\Gamma_i^{H \rightarrow X}$ ,  $\Gamma_i^H$  are estimated using Monte Carlo simulations in the next section.

## 12.4 Interpretation within the SM-EFT

The statistical model used to obtain the signal strength measurements in the STXS bins, depicted in Fig. 12.2, is reinterpreted in terms of the SM-EFT and used to constrain the Wilson coefficients of SM-EFT operators. Using the linearized SM-EFT model described in § 12.3, the signal strength in each STXS category is reinterpreted in terms of deviations from the SM value as a function of the SM-EFT operators that impact the signal cross-section to determine its magnitude.

### 12.4.1 Simulating the impact of dimension-6 SM-EFT operators

All operators included in the fit can be found in Tab. 12.2. The operators in this mea-

**Table 12.2:** The Wilson coefficients ( $c_i$ ) included in the SM-EFT interpretation of the Higgs combination and their corresponding dimension-6 SM-EFT operators ( $\mathcal{O}_i^{(6)}$ ).

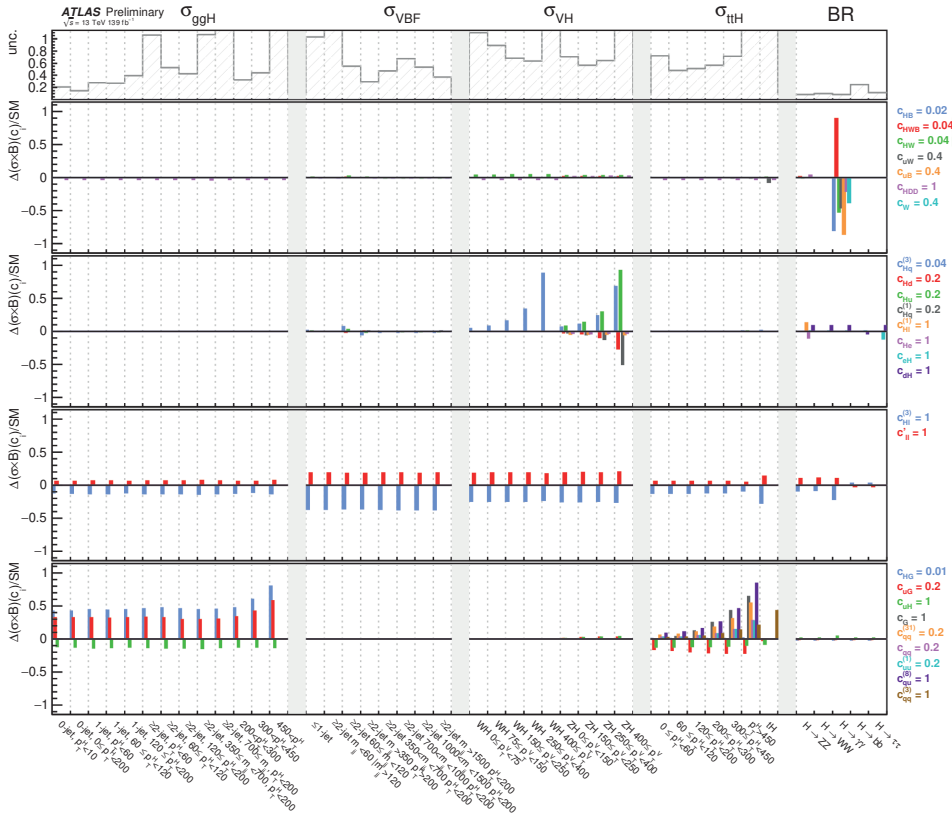
| $c_i$          | Operator $\mathcal{O}_i^{(6)}$                                               | $c_i$           | Operator $\mathcal{O}_i^{(6)}$                                   |
|----------------|------------------------------------------------------------------------------|-----------------|------------------------------------------------------------------|
| $c_{H\Box}$    | $(H^\dagger H)\Box(H^\dagger H)$                                             | $c_{uG}$        | $(\bar{q}_p\sigma^{\mu\nu}T^A u_r)\tilde{H}G_{\mu\nu}^A$         |
| $c_{HDD}$      | $(H^\dagger D^\mu H)^*(H^\dagger D_\mu H)$                                   | $c_{uW}$        | $(\bar{q}_p\sigma^{\mu\nu}u_r)\tau^I\tilde{H}W_{\mu\nu}^I$       |
| $c_{HG}$       | $H^\dagger H G_{\mu\nu}^A G^{A\mu\nu}$                                       | $c_{uB}$        | $(\bar{q}_p\sigma^{\mu\nu}u_r)\tilde{H}B_{\mu\nu}$               |
| $c_{HB}$       | $H^\dagger H B_{\mu\nu}B^{\mu\nu}$                                           | $c_{ll}^{(1)}$  | $(\bar{l}_p\gamma_\mu l_t)(\bar{l}_r\gamma^\mu l_s)$             |
| $c_{HW}$       | $H^\dagger H W_{\mu\nu}^I W^{I\mu\nu}$                                       | $c_{qq}^{(1)}$  | $(\bar{q}_p\gamma_\mu q_t)(\bar{q}_r\gamma^\mu q_s)$             |
| $c_{HWB}$      | $H^\dagger \tau^I H W_{\mu\nu}^I B^{\mu\nu}$                                 | $c_{qq}^{(3)}$  | $(\bar{q}_p\gamma_\mu\tau^I q_r)(\bar{q}_s\gamma^\mu\tau^I q_t)$ |
| $c_{eH}$       | $(H^\dagger H)(\bar{l}_p e_r \tilde{H})$                                     | $c_{qq}$        | $(\bar{q}_p\gamma_\mu q_t)(\bar{q}_r\gamma^\mu q_s)$             |
| $c_{uH}$       | $(H^\dagger H)(\bar{q}_p u_r \tilde{H})$                                     | $c_{qq}^{(31)}$ | $(\bar{q}_p\gamma_\mu\tau^I q_t)(\bar{q}_r\gamma^\mu\tau^I q_s)$ |
| $c_{dH}$       | $(H^\dagger H)(\bar{q}_p d_r \tilde{H})$                                     | $c_{uu}$        | $(\bar{u}_p\gamma_\mu u_r)(\bar{u}_s\gamma^\mu u_t)$             |
| $c_{Hl}^{(1)}$ | $(H^\dagger i\overleftrightarrow{D}_\mu H)(\bar{l}_p\gamma^\mu l_r)$         | $c_{uu}^{(1)}$  | $(\bar{u}_p\gamma_\mu u_t)(\bar{u}_r\gamma^\mu u_s)$             |
| $c_{Hl}^{(3)}$ | $(H^\dagger i\overleftrightarrow{D}_\mu^I H)(\bar{l}_p\tau^I\gamma^\mu l_r)$ | $c_{qu}^{(1)}$  | $(\bar{q}_p\gamma_\mu q_t)(\bar{u}_r\gamma^\mu u_s)$             |
| $c_{He}$       | $(H^\dagger i\overleftrightarrow{D}_\mu H)(\bar{e}_p\gamma^\mu e_r)$         | $c_{ud}^{(8)}$  | $(\bar{u}_p\gamma_\mu T^A u_r)(\bar{d}_s\gamma^\mu T^A d_t)$     |
| $c_{Hq}^{(1)}$ | $(H^\dagger i\overleftrightarrow{D}_\mu H)(\bar{q}_p\gamma^\mu q_r)$         | $c_{qu}^{(8)}$  | $(\bar{q}_p\gamma_\mu T^A q_r)(\bar{u}_s\gamma^\mu T^A u_t)$     |
| $c_{Hq}^{(3)}$ | $(H^\dagger i\overleftrightarrow{D}_\mu^I H)(\bar{q}_p\tau^I\gamma^\mu q_r)$ | $c_{qd}^{(8)}$  | $(\bar{q}_p\gamma_\mu T^A q_r)(\bar{d}_s\gamma^\mu T^A d_t)$     |
| $c_{Hu}$       | $(H^\dagger i\overleftrightarrow{D}_\mu H)(\bar{u}_p\gamma^\mu u_r)$         | $c_W$           | $\epsilon^{JK}W_\mu^{I\nu}W_\nu^{J\rho}W_\rho^{K\mu}$            |
| $c_{Hd}$       | $(H^\dagger i\overleftrightarrow{D}_\mu H)(\bar{d}_p\gamma^\mu d_r)$         | $c_G$           | $f^{ABC}G_\mu^{A\nu}G_\nu^{B\rho}G_\rho^{C\mu}$                  |

surement are defined using the Warsaw Basis for CP-even mass dimension-6 operators [250] and constitute operators for which the absolute  $O(\Lambda^{-2})$  impact of a unit deviation of the coefficient is at least 0.1% of the total SM cross-section in any STXS region. In total, 64 operators are found to either directly or indirectly impact the Higgs boson production or decay within these criteria. The operators' impacts have been computed

using MadGraph5\_aMC@NLO version 2.8.1 [109], interfaced to PYTHIA 8 [110], with the A14 tune [113] for parton showering and hadronization.

Calculations for the Higgs processes that have tree level diagrams for QCD and for EW interactions are calculated using SMEFTsimv3.0 [251] using the  $U(3)^5$  flavour symmetry. Processes that do not occur on tree level, like  $ggF$ ,  $gg \rightarrow ZH$  and  $H \rightarrow gg$  are performed at NLO accuracy in QCD using SMEFT@NLO[252]. The  $ggF$  process is generated with zero, one, and two additional partons in the final state. The MLM merging algorithm [253, 254] is used to eliminate overlap between the matrix elements.

Finally, the calculations of SM-EFT and SM interference for  $H \rightarrow \gamma\gamma$  are performed at NLO accuracy in QED [255]. Contributions of SM-EFT to the background processes are not considered. The impact of each operator on the signal production cross-section, total width  $\Gamma_H^{tot}$ , partial width  $\Gamma_{H \rightarrow X}$  and eventual acceptance corrections in the STXS stage-1.2 categories are shown in Fig. 12.3.



**Figure 12.3:** An illustration of the impact of the most relevant SM-EFT operators on the STXS regions and decay modes, relative to the SM cross-section. In order to judge the experimental sensitivity available to constrain the operators from the data in the listed STXS regions, the total uncertainty on measurement in the corresponding regions is shown in the top panel.



The parton-level events generated used to assess the operator impact are not simulated and reconstructed using ATLAS GEANT4 reconstruction software. Instead, they are studied at particle-level after being passed through the shower and hadronization in PYTHIA 8. This is possible, as STXS signal strengths represent particle-level information, which is unfolded from the data through a fit of multiple particle-level contributions to reconstruction-level regions. The event categorisation into the various STXS categories is performed using a RIVET analysis [256]. The RIVET analysis is based on the ATLAS analysis code and is used to categorize the events using particle-level information of the Higgs boson and its associated particles in the production. The acceptance of the STXS categories is calculated under the SM hypothesis, unless otherwise specified.

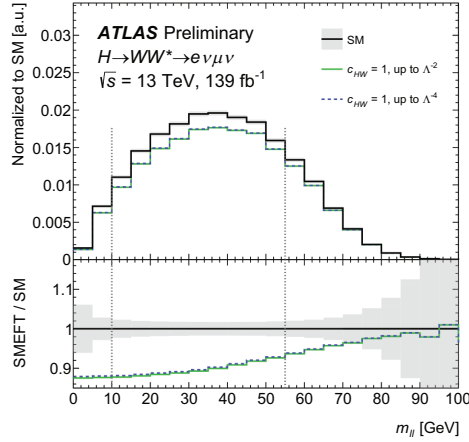
#### 12.4.2 Acceptance effects in $H \rightarrow WW^*$ and $H \rightarrow ZZ^*$ measurements

The STXS binning of the Higgs production is defined such that the acceptance of targeted processes can be assumed to be flat in each STXS category. Therefore, the acceptance effects computed in each analysis, included in the combination, are calculated under the SM hypothesis. Within the framework of SM-EFT, these acceptance corrections are also assumed to cover potential acceptance differences due to contributions of BSM physics to Higgs production. Conversely, the STXS framework cannot model acceptance effects in Higgs decays, as it does not define fiducial regions in the decay phase-space. Some analyses in the combination target very specific final state signatures based on the SM expectation of these decay kinematics.

For most decays, in particular the 2-body decays such as  $H \rightarrow b\bar{b}$ ,  $H \rightarrow c\bar{c}$  or  $H \rightarrow \gamma\gamma$ , acceptance effects are found to be negligible. However, the acceptance of  $H \rightarrow WW^*$  and  $H \rightarrow ZZ^*$  decays changes significantly for a small set of SM-EFT operators. This sensitivity arises because the analysis selections are tailored towards the SM  $H \rightarrow WW^*$  and  $H \rightarrow ZZ^*$  decay topology. Some SM-EFT operators introduce point-like interactions between Higgs bosons and leptons and can cause potentially different lepton kinematics. The  $H \rightarrow WW^*$  selection criteria on the di-lepton invariant mass of  $10 \text{ GeV} < m_{\ell\ell} < 55 \text{ GeV}$ , which is used to select the on-shell Higgs peak and cut out large amounts of  $t\bar{t}$  background events, has a significant effect on the acceptance for operator  $\mathcal{O}_{HW}$ . The corresponding differential distribution in  $m_{ll}$  is shown in Fig. 12.4.

The other operators affecting the  $H \rightarrow WW^*$  decay,  $\mathcal{O}_{HDD}$ ,  $\mathcal{O}_u^{(1)}$ ,  $\mathcal{O}_{H\Box}$ , change the overall rate of the Higgs production uniformly throughout the phase-space and therefore do not result in any additional acceptance effects. The  $H \rightarrow ZZ^*$  analysis has stringent requirements on the mass of the reconstructed  $Z$  bosons, amounting to  $m_{Z_2} > 12 \text{ GeV}$ , which severely decreases the impact of the  $\mathcal{O}_{HW}$  operator. Because of these effects, the  $H \rightarrow WW^*$  and  $H \rightarrow ZZ^*$  analyses performed dedicated acceptances studies resulting into a set of acceptance corrections for both of the decay channels.

The corrections are calculated by estimating the operator impact inclusively and after applying the analysis selections on particle-level, and define the difference as the acceptance



**Figure 12.4:** Illustration of the order- $\Lambda^2$  and  $\Lambda^4$  effects of the  $\mathcal{O}_{HW}$  operator on the shape of kinematic observables used in the  $H \rightarrow WW^*$  analysis. The window selection, made using the di-lepton invariant mass in the  $H \rightarrow WW^*$  measurement, is also shown. The grey band represents the statistical uncertainty of the SM sample. Taken from Ref. [257].

correction. The comparison of the operator impact with and without acceptance corrections up to order  $\Lambda^{-2}$ , inclusively for  $H \rightarrow ZZ^*$  and per STXS region for  $H \rightarrow WW^*$ , is shown in Fig. 12.5.

The parameterization of  $\sigma_{\text{SM-EFT}}^p \cdot \mathcal{B}_{\text{SM-EFT}}^{H \rightarrow X}$  can be rewritten into the total number of events in an STXS region,  $N_{\text{SM-EFT}}^{\text{STXS}}$  and total operator impact,  $A_i$ , on this yield

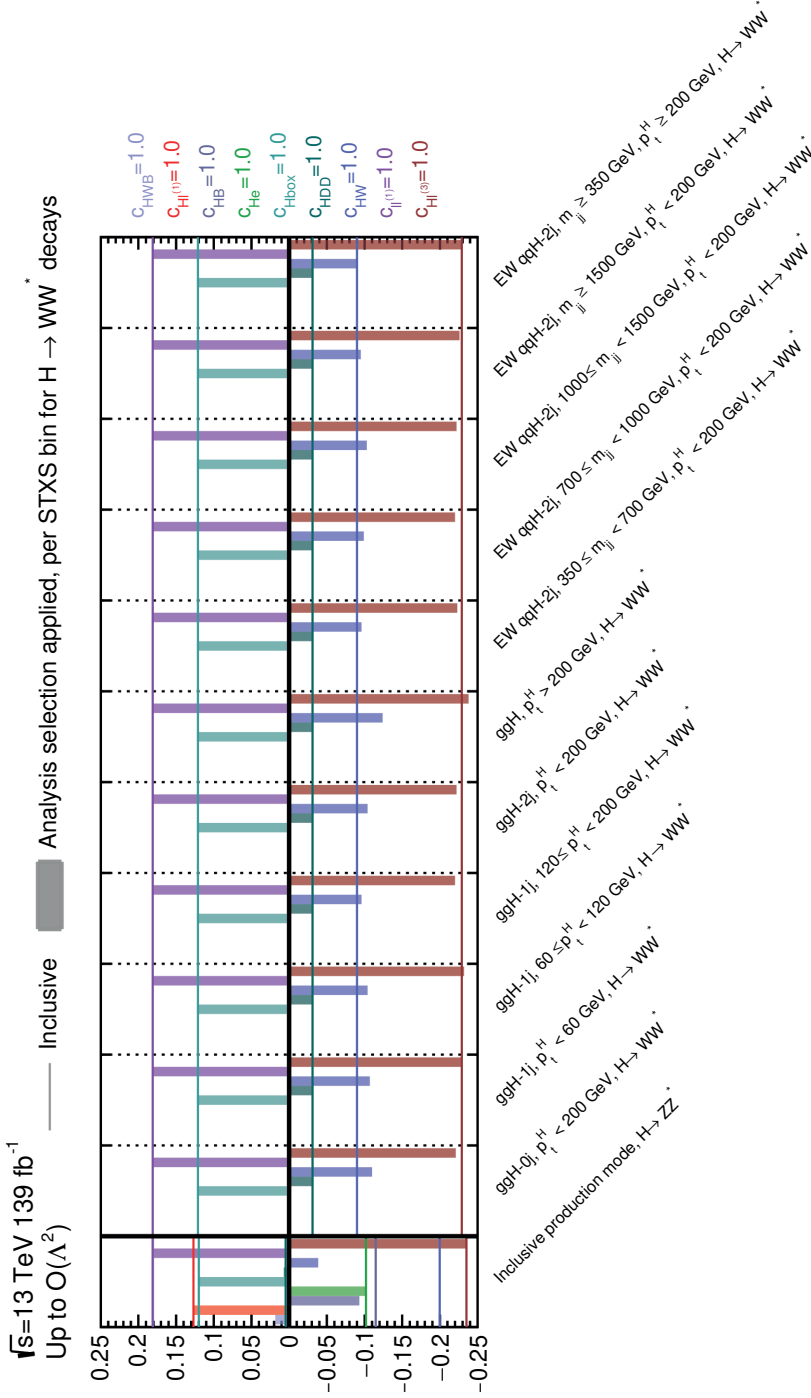
$$N_{\text{SM-EFT}}^{\text{STXS}} = N_{\text{SM}}^{\text{STXS}} \cdot \left( 1 + \sum_i A_i^{\text{STXS}} c_i + \sum_i (A_i^{H \rightarrow X} - A_i^H) \cdot c_i \right),$$

and becomes

$$N_{\text{SM-EFT}}^{\text{STXS}} = N_{\text{SM}}^{\text{STXS}} \cdot \left( 1 + \sum_i A_i^{\text{STXS}} c_i + \sum_i (A_i^{H \rightarrow X} - A_i^H - A_i^{\text{acceptance}}) \cdot c_i \right),$$

when the acceptance correction after linearization is added. The maximal acceptance effect on the operator impact amounts to approximately 10% for  $H \rightarrow WW^*$  and 75% for  $H \rightarrow ZZ^*$ .

The acceptance correction in the profile likelihood fit is applied separately for each Higgs production process, as the large differences in the reconstruction of the ggF and VBF processes results in sizeably different acceptance effects. The acceptance effects on the  $H \rightarrow ZZ^*$  measurement are found to be uniform across the STXS regions and are therefore applied inclusively. These corrections are applied as the linear corrections to the decay parameterization and are shown in Tab. 12.3.



**Figure 12.5:** Impact of the SM-EFT operators on the partial width  $\Gamma^{H \rightarrow X^*}$ , in the  $H \rightarrow WW^*$  and  $H \rightarrow ZZ^*$  measurements. The impact is shown relative to the SM prediction. In this figure, the acceptance correction is evaluated separately per STXS region for the  $H \rightarrow WW^*$  measurement and inclusively for the  $H \rightarrow ZZ^*$  measurement. The horizontal lines represent the operator impact without acceptance corrections, while coloured bars show the impact with acceptance corrections.

**Table 12.3:** Parameterization of the acceptance corrections in the relevant Higgs boson decay channels in terms of the EFT parameters defined in the SMEFTsimmodel. Corrections are obtained as a difference of the fiducial and the inclusive signal predictions for the considered decay vertex using the fully linearized signal parameterization approach.

| Decay mode                                      | Acceptance parameterization                |
|-------------------------------------------------|--------------------------------------------|
| $H \rightarrow 4l$                              | $0.162c_{HW} + 0.021c_{HB} - 0.162c_{HWB}$ |
| $ggH, H \rightarrow WW^* \rightarrow l\nu l\nu$ | $-0.007c_{HW} + 0.006c_{Hl}^{(3)}$         |
| $VBF, H \rightarrow WW^* \rightarrow l\nu l\nu$ | $-0.018c_{HW} + 0.007c_{Hl}^{(3)}$         |

### 12.4.3 The statistical model of the STXS SM-EFT interpretation

The statistical model of the STXS SM-EFT measurement is a reparameterized version of the likelihood function of the STXS cross-section measurement. For each analysis bin  $b$  of a STXS signal, the likelihood associated with the observation of  $N_b$  signal events is modelled as a Poisson distribution

$$L_b(N_b | \{\mu_p\}, \theta^{\text{STXS}}) = \text{Poisson} \left( N_b \left| \sum_p \mu_p S_{p,b}(\theta^{\text{STXS}}) + B_b(\theta^{\text{STXS}}) \right. \right),$$

where  $B_b$  is the expected background yield and  $S_{p,b}$  the expected number of Higgs signal events, where the index  $p$  runs over all Higgs production processes (ggF, VBF,  $VH$ ,  $t\bar{t}H$  or  $tH$ ). The parameter  $\mu_p$  is the signal strength parameter and represents the cross-section times branching ratio of the production process divided by the SM expectation. Both  $B_b$  and  $S_{p,b}$  depend on a set of nuisance parameters  $\theta^{\text{STXS}}$ , as both the expected signal and background yield are subject to uncertainties in the modelling of experimental effects and in the theoretical prediction of signal, as well as background, processes.

For the purpose of the SM-EFT interpretations, the signal strength parameters  $\mu_p$  is reparameterized in terms of Wilson coefficients,

$$\mu_p(\mathbf{c}) = 1 + \sum_i A_i^p c_i + \sum_i c_i \Gamma_i^{H \rightarrow X} - \sum_i c_i \Gamma_i^{\text{tot}} - \sum_i A_i^{\text{acceptance}} c_i,$$

here the index  $i$  runs over all considered SM-EFT operators. The joint likelihood for all STXS regions, which includes all  $n_{\text{bins}}^{\text{SR+CR}}$  analysis bins in signal and control regions as well as constraint terms  $f_i$  on nuisance parameters, then becomes

$$L(\mathbf{N} | \mathbf{c}, \theta^{\text{STXS}}) = \prod_b^{n_{\text{bins}}^{\text{SR+CR}}} \text{Poisson} \left( N_b \left| \sum_p \mu_p(\mathbf{c}) S_{p,b}(\theta^{\text{STXS}}) + B_b(\theta^{\text{STXS}}) \right. \right) \times \prod_i^{n_{\text{syst}}^{\text{STXS}}} f_i(\theta_i^{\text{STXS}}),$$

where  $\mathbf{N} = (N_1, \dots, N_b, \dots, N_{n_{\text{bins}}^{\text{SR+CR}}})$  is the observed number of events in the analysis bins. The resulting likelihood,  $L(\mathbf{N} | \mathbf{c}, \theta^{\text{STXS}})$ , is then expressed in Wilson coefficients.

The estimation of the confidence intervals for a Wilson coefficient  $c_i$  is done with the use of a profile likelihood ratio test statistic constructed from the likelihood:

$$\lambda(c_i) = -2 \log \frac{L(c_i, \hat{\hat{\theta}}(c_i))}{L(\hat{c}_i, \hat{\theta}(c_i))}.$$

Here  $L(c_i, \hat{\hat{\theta}})$  is the maximum of the likelihood for a chosen  $c_i$  and  $L(\hat{c}_i, \hat{\theta})$  the value at the absolute maximum of the likelihood. Following Wilks' theorem, confidence intervals at 68% (95%) on parameters  $c_i$  are defined as the intervals in the parameter of  $c_i$  that correspond to a rise of at most 0.5 (1.96) units in the the profile likelihood ratio test statistic  $\lambda(c_i)$ . For intervals on a subset of  $c_i$ , e.g. one-dimensional intervals on a single  $c_i$ , the other Wilson coefficients are treated as nuisance parameters.

## 12.5 Systematics

Systematics uncertainties that affect measurements across multiple STXS regions are modelled with common nuisance parameters and originate from the reconstruction algorithms, luminosity measurements, object energy scale or identification efficiencies. The correlated effect of background systematics across different channels is evaluated on a case by case basis for both experimental and theoretical uncertainties.

## 12.6 Eigenvector decomposition

Because the available data in the analyses contain insufficient information to simultaneously constrain all coefficients  $c_i$  of the operators in [Tab. 12.2](#), an eigenvector (EV) decomposition of the SM-EFT operator basis is performed. As Warsaw basis operators, which are defined by theoretical considerations, generally align poorly with experimental observables, it is common that individual operator coefficients, as well as linear combinations of coefficients, are unconstrained by a given set of measurements. As the existence of unconstrained (i.e. flat) directions in the parameter space cause numerical likelihood maximization algorithms to fail, these must be explicitly removed prior to parameter estimation. In this section, the methodology to identify and remove these are explained.

First, an eigenvalue decomposition is performed on the SM expected covariance matrix of the cross-section STXS measurement  $(\sigma \times B)_{if}$ , assuming Gaussian behaviour. The eigenvectors of the Fisher information matrix provide a list of orthogonal directions in the Wilson coefficient parameters ranked by experimental sensitivity, as the eigenvalues are proportional to the inverse of the expected measurement variance in the corresponding direction. The covariance matrix  $V_{\text{measurement}}$  of the Higgs signal strength parameters is inverted to form the Hessian matrix  $V_{\text{measurement}}^{-1}$ . This matrix is then reparameterized in terms of Wilson coefficients using the parameterization matrix  $P_{\mu^x \rightarrow c_i}$ , which maps the signal strength parameters to linear combinations of EFT coefficients to obtain the Fisher information matrix

$$V_{\text{SM-EFT}}^{-1} = P_{\mu^x \rightarrow c_i}^T V_{\text{measurement}}^{-1} P_{\mu^x \rightarrow c_i}.$$

In order to reduce the impact of unwanted correlations caused by operators with low sensitivity, parameter groups are formed of similar physical meaning.

The  $c_{H\Box}$  operator modifies the global Higgs boson production cross-section, constituting an overall normalization effect to all STXS regions, and can be added back to the fit as a correction to all parameters. Although it presents a degeneracy in the fit, it does not affect the final fit basis.

The  $c_{Hq}^{(3)}$  operator represents a point interaction between the Higgs boson and the quarks, showing strong sensitivity in the  $H \rightarrow b\bar{b}$  VH measurements and is therefore not grouped with other parameters in the fit.

The  $c_{eH}$  operator is mainly constrained by the  $H \rightarrow \tau\tau$  measurements, which are sensitive to the Yukawa coupling  $Y_\tau$ . Due to the lack of sensitivity of  $c_{eH}$  in other measurements, it is kept as separate parameter in the fit model. Similarly, the  $c_{dH}$  operator is constrained by the measurement of the bottom Yukawa coupling in the  $H \rightarrow b\bar{b}$  measurements and is kept as a separate parameter in the fit model.

The operators  $c_{HG}$ ,  $c_{uG}$  and  $c_{uH}$  are constrained by the ggF and some  $t\bar{t}H$  measurements, showing a strong dependence on the  $p_T^H$  of the probed STXS regions. These operators are grouped together into a parameter-subgroup in the fit model, with most of their sensitivity found in analyses providing measurements of the ggF cross-section in the 0 and 1-jet bin, primarily the  $H \rightarrow WW^*$  and  $H \rightarrow \gamma\gamma$  measurements.

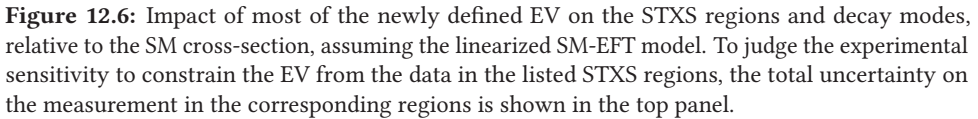
The set  $c_{qq}^{(1)}$ ,  $c_{qq}$ ,  $c_{qq}^{(3)}$ ,  $c_{qq}^{(31)}$ ,  $c_{uu}$ ,  $c_{uu}^{(1)}$ ,  $c_{ud}^{(8)}$ ,  $c_{qu}^{(1)}$ ,  $c_{qu}^{(8)}$ ,  $c_{qd}^{(8)}$  and  $c_G$  is constrained by the  $t\bar{t}H$  measurements, affecting the shape of the signal similarly. The updated  $t\bar{t}H$ ,  $H \rightarrow b\bar{b}$  measurements introduce additional separation power between the ggF and  $t\bar{t}H$  related parameters. These operators are grouped together in the  $c_{top}$  group and the most sensitive eigenvector of this parameter group is introduced as parameter in the fit model.

The  $c_{uW}$  and  $c_{uB}$  operators are constrained via the  $t\bar{t}H$  and  $H \rightarrow \gamma\gamma$  channels. Besides  $c_{uW}$  and  $c_{uB}$ , the  $H \rightarrow \gamma\gamma$  channel also constrains the  $c_{HW}$ ,  $c_{HB}$ ,  $c_{HWB}$ ,  $c_W$  and  $c_{HDD}$ , leading to a joint rotation of these six Wilson coefficients. These operators affect the  $H \rightarrow \gamma\gamma$  branching ratio and the VBF and  $VH$  production modes, well measured in the  $H \rightarrow WW^*$  and  $H \rightarrow b\bar{b}$  channels respectively. The three most sensitive eigenvectors of this parameter group are introduced as parameters in the fit model.

Operator set  $c_{Hl}^{(1)}$ ,  $c_{He}$ ,  $c_{Hu}$ ,  $c_{Hd}$ , and  $c_{Hq}^{(1)}$  is mainly constrained by the  $VH$ ,  $H \rightarrow b\bar{b}$  and VBF  $H \rightarrow b\bar{b}$  measurements, where strong correlations exist between the  $c_{Hu}$ ,  $c_{Hd}$  and  $c_{Hq}^{(1)}$  parameters. The first two operators introduce modifications to the Yukawa couplings  $Y_{q=u,c,t}$  and  $Y_{d=s,b}$ , while  $c_{Hq}^{(1)}$  introduces a 4-point interaction between a Higgs,  $W$  boson, and a top and bottom quark. Thus, these operators are grouped into two sets: one affecting quarks  $\{c_{Hu}, c_{Hd}, c_{Hq}^{(1)}\}$  and the other leptons  $\{c_{Hl}^{(1)}, c_{He}\}$ .

Finally,  $c_{Hl}^{(3)}$  and  $c_{ll}^{(1)}$  are grouped together because they both introduce a shift in the Fermi constant  $\Delta G_F$ , with  $c_{Hl}^{(3)}$  representing a positive and  $c_{ll}^{(1)}$  a negative shift, making

The impact of the new fit basis on the STXS measurement, shown in Fig. 12.6, is more aligned with the STXS signal strengths, their respective sensitivities and parameters and can therefore be constrained with only small to moderate correlations between the estimates. The list of eigenvectors is truncated at  $\lambda = 0.1$ , discarding all directions with an expected experimental uncertainty larger than 10. The kept eigenvectors are shown in Fig. 12.7, while the absence of sensitivity in the discarded directions is explicitly verified with a likelihood scan. After the rotation into the new basis, the fit is performed using the new fit basis consisting of three Wilson coefficients and 10 subgroup eigenvectors.



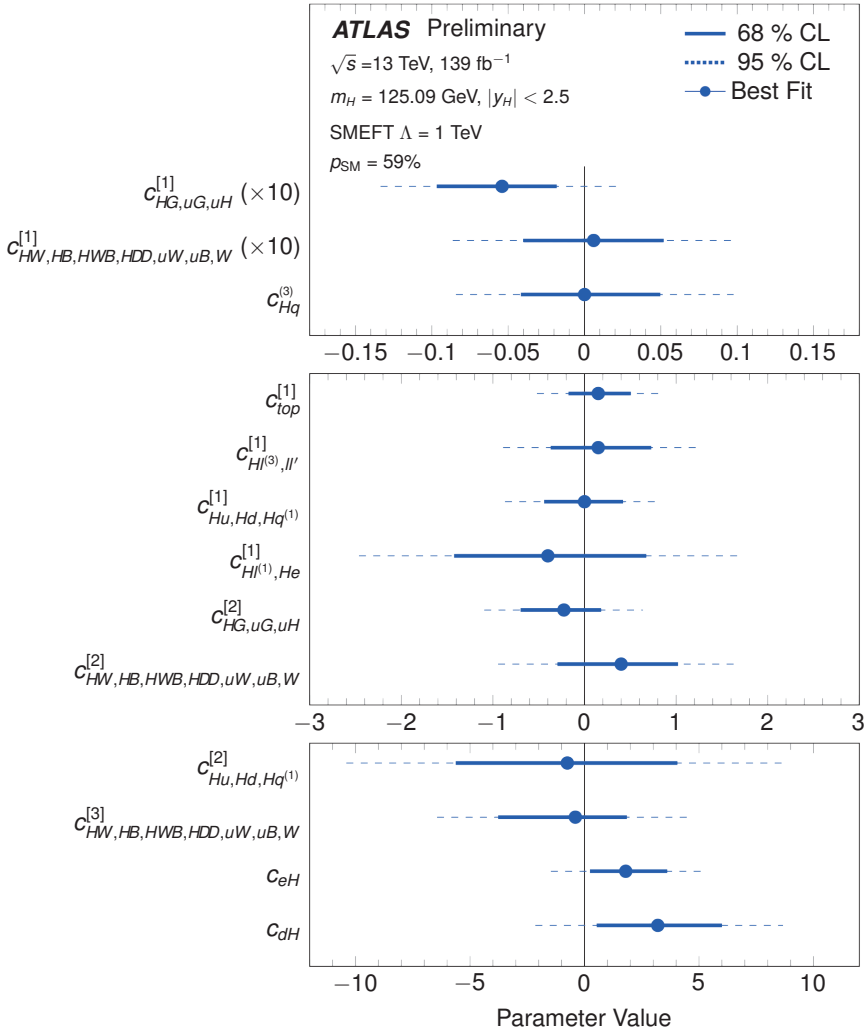
**Figure 12.6:** Impact of most of the newly defined EV on the STXS regions and decay modes, relative to the SM cross-section, assuming the linearized SM-EFT model. To judge the experimental sensitivity to constrain the EV from the data in the listed STXS regions, the total uncertainty on the measurement in the corresponding regions is shown in the top panel.





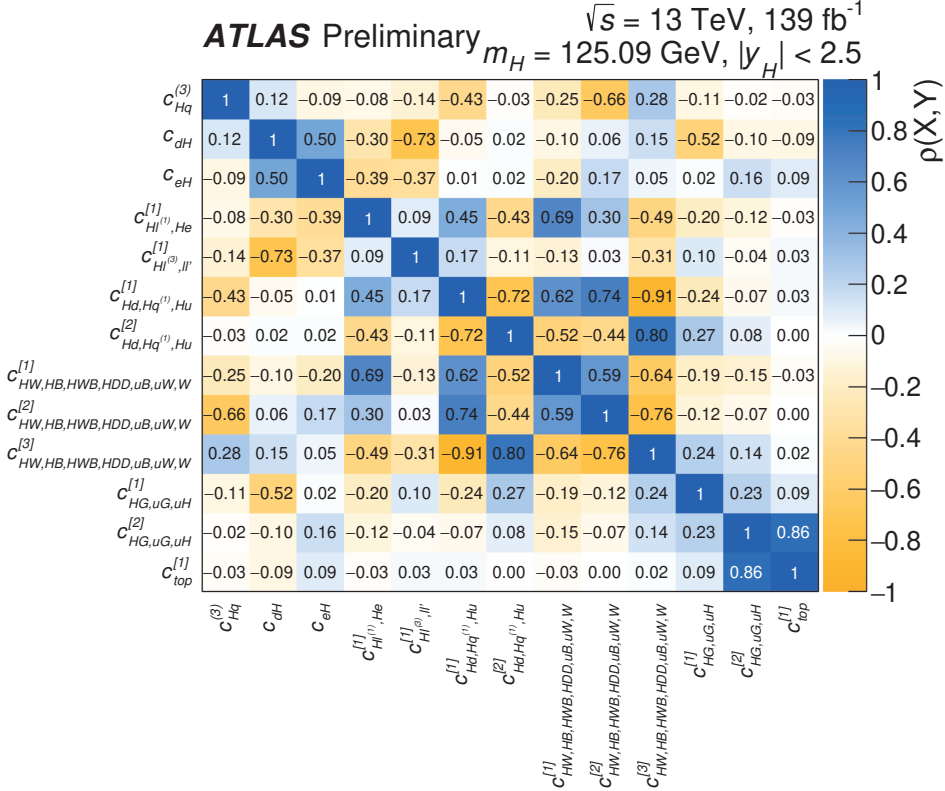
## 12.7 SM-EFT interpretation results

A simultaneous fit using the 13 parameters of the fit basis is performed to the data. The results of the fit, along with the correlations between the operators,  $\rho(c_i, c_j)$ , are shown in Figs. 12.8 and 12.9, as well as in Tab. 12.4. The results are obtained by performing one-dimensional scans of the profile likelihood ratio,  $\lambda(c_i)$ , now as a function of the newly defined EV directions, and are presented in the form of 68% and 95% confidence intervals.



**Figure 12.8:** Summary of observed measurements of the parameters  $c'_i$  within the SM-EFT linearized model. The ranges shown correspond to 68% (solid) and 95% (dashed) confidence level intervals.

In each of the profile likelihood scans, the nuisance parameters as well as all other 12 parameters of interest are profiled. The level of compatibility between the SM hypothesis and the best-fit point corresponds to a  $p$ -value of  $p_{\text{SM}} = 59\%$ , using the 13 degrees of freedom as defined in the previous section. The SM prediction, for all EV directions, is found within the 95% confidence interval.



**Figure 12.9:** The correlation coefficients in  $\rho(X, Y)$  from the linearized SM-EFT model between observables for the observed data.

**Table 12.4:** Summary of the observed and the SM expected measurement of the parameters  $c'_i$  with the SM-EFT linearized model. Showing 68% and 95% confidence level intervals.

| Model Parameter                   |          | Observed           |                    | Expected          |                   |
|-----------------------------------|----------|--------------------|--------------------|-------------------|-------------------|
| $\Lambda = 1 \text{ TeV}$         | Best-fit | 68% CI             | 95% CI             | 68% CI            | 95% CI            |
| $c_{Hq}^{(3)}$                    | 0.0      | $[-0.04, 0.05]$    | $[-0.08, 0.1]$     | $[-0.04, 0.05]$   | $[-0.08, 0.09]$   |
| $c_{dH}$                          | 3.2      | $[0.5, 6]$         | $[-2.1, 9]$        | $[-2.7, 2.7]$     | $[-5, 5]$         |
| $c_{eH}$                          | 1.8      | $[0.23, 4]$        | $[-1.5, 5]$        | $[-1.7, 1.7]$     | $[-3.5, 3.2]$     |
| $c_{HW,HB,HWB,HDD,uW,uB,W}^{[1]}$ | 0.001    | $[-0.004, 0.005]$  | $[-0.009, 0.01]$   | $[-0.005, 0.004]$ | $[-0.009, 0.009]$ |
| $c_{HW,HB,HWB,HDD,uW,uB,W}^{[2]}$ | 0.4      | $[-0.30, 1.0]$     | $[-0.9, 1.7]$      | $[-0.6, 0.6]$     | $[-1.3, 1.3]$     |
| $c_{HW,HB,HWB,HDD,uW,uB,W}^{[3]}$ | -0.4     | $[-4, 1.9]$        | $[-6, 5]$          | $[-2.7, 2.8]$     | $[-5, 6]$         |
| $c_{Hl^{(0)},He}^{[1]}$           | -0.4     | $[-1.4, 0.7]$      | $[-2.5, 1.7]$      | $[-1.0, 1.0]$     | $[-2.0, 2.0]$     |
| $c_{Hu,Hd,Hq^{(1)}}^{[1]}$        | 0.0      | $[-0.4, 0.4]$      | $[-0.9, 0.8]$      | $[-0.4, 0.4]$     | $[-0.9, 0.8]$     |
| $c_{Hu,Hd,Hq^{(2)}}^{[2]}$        | -0.8     | $[-6, 4]$          | $[-10, 9]$         | $[-5, 5]$         | $[-10, 10]$       |
| $c_{Hl^{(3)},ll'}^{[1]}$          | 0.15     | $[-0.4, 0.7]$      | $[-0.9, 1.3]$      | $[-0.5, 0.5]$     | $[-1.0, 1.0]$     |
| $c_{HG,uG,uH}^{[1]}$              | -0.005   | $[-0.01, -0.0018]$ | $[-0.013, 0.0021]$ | $[-0.004, 0.004]$ | $[-0.008, 0.008]$ |
| $c_{HG,uG,uH}^{[2]}$              | -0.23    | $[-0.7, 0.18]$     | $[-1.1, 0.6]$      | $[-0.4, 0.5]$     | $[-0.9, 0.9]$     |
| $c_{top}^{[1]}$                   | 0.15     | $[-0.18, 0.5]$     | $[-0.5, 0.8]$      | $[-0.4, 0.4]$     | $[-0.7, 0.7]$     |

## COMBINED SM-EFT INTERPRETATION OF HIGGS AND $WW$ MEASUREMENTS

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This chapter presents the first-ever Run-2 EW + Higgs effective field theory interpretation, in which Wilson coefficients are constrained from two different analyses with dissimilar analysis techniques [8]. This is achieved by combining a measurement of  $H \rightarrow WW^* \rightarrow e^\pm \nu_e \mu^\mp \nu_\mu$  in the gluon-fusion and vector-boson fusion production modes with differential cross-section measurements of  $WW$  diboson production in the  $e^\pm \nu_e \mu^\mp \nu_\mu$  decay channel, with a veto on additional jets.

### 13.1 Data and input measurements

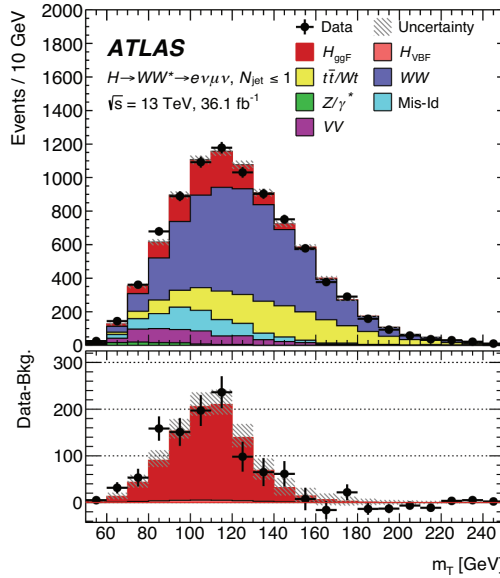
The combination uses measurements of an on-shell Higgs boson ( $H$ ) decaying into a pair of  $W$  bosons ( $H \rightarrow WW^*$ ) [258] and of two on-shell  $W$  bosons ( $WW$ ) [259]. These measurements are based on  $36.1 \text{ fb}^{-1}$  of  $pp$ -collision data collected at a centre-of-mass energy of  $\sqrt{s} = 13 \text{ TeV}$  by the ATLAS experiment at the LHC during 2015 and 2016. Combining these measurements presents several challenges: overlapping event selections,  $WW$  background to  $H \rightarrow WW^*$  affecting EFT interpretations, correlated systematic uncertainties, and the need to integrate different analysis techniques into a coherent statistical framework for SM-EFT interpretations.

#### 13.1.1 The $H \rightarrow WW^*$ couplings measurement

The  $H \rightarrow WW^*$  measurement targets Higgs boson production via ggF and VBF, with decays through  $H \rightarrow WW^* \rightarrow e^\pm \nu_e \mu^\mp \nu_\mu$ . It is performed in three signal regions with zero, one, and two jets. Signal candidates are selected by requiring an isolated  $e^\pm \mu^\mp$  pair, with transverse momentum thresholds of 22 and 15 GeV for the leading and subleading-leptons, respectively. Events with jets tagged as containing  $b$ -hadrons are rejected to suppress top quark background contributions and various selections designed to enhance the  $H \rightarrow WW^*$  topology are made. The data to MC comparison in the  $N_{\text{jet}} \leq 1$  ggF signal region is shown in Fig. 13.1.

#### 13.1.2 The $WW$ differential measurement

The  $WW$  measurement accepts all production modes decaying into  $e^\pm \nu_e \mu^\mp \nu_\mu$  final states. To suppress  $H \rightarrow WW^*$  production, the invariant mass of the two charged leptons,  $m_{e\mu}$ , is required to be greater than 55 GeV, ensuring the production of two on-shell  $W$  bosons. A jet veto with  $p_T > 35 \text{ GeV}$  and  $|\eta| < 4.5$  is applied to suppress



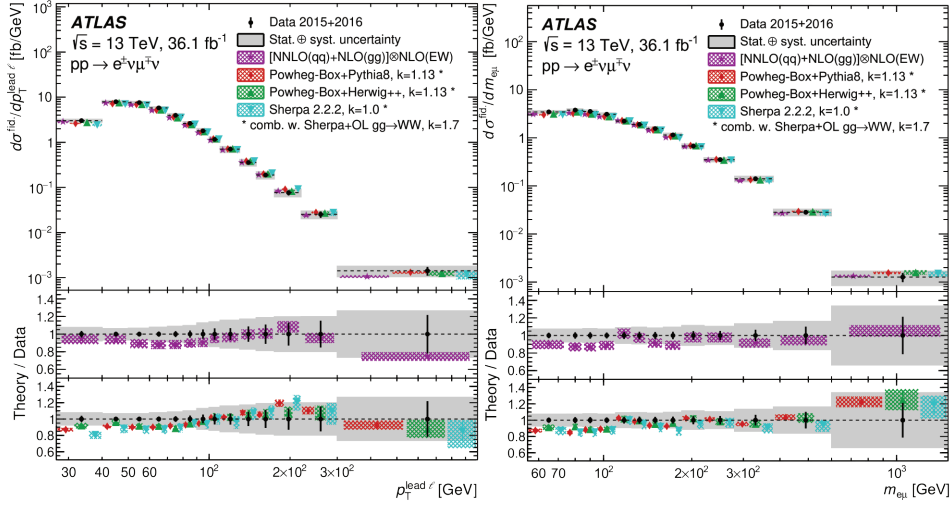
**Figure 13.1:** The post-fit transverse mass  $m_T$  distribution for the combined  $N_{\text{jet}} \leq 1$  ggF signal regions. The bottom panel shows the data subtracted background modelling for the distribution, assuming a SM Higgs boson with  $m_H = 125$  GeV. The hatched band illustrates the uncertainty on signal and background [258].

top quark backgrounds. After background subtraction, the data is unfolded and fiducial cross-sections are reported for a particle-level phase-space similar to the reconstruction level selection. The unfolded  $p_T^{\text{lead. lep.}}$  and  $m_{e\mu}$  distributions are shown in Fig. 13.2.

In this reinterpretation, the unfolded differential cross-section as a function of the highest lepton transverse momentum,  $p_T^{\text{lead. lep.}}$ , is used in the fit model. This distribution is the most sensitive to dimension-6 operator effects that increase with the parton centre-of-mass energy. The differential cross-section is measured in 14 bins, with the lowest threshold at 27 GeV and with the last bin including all events with  $p_T^{\text{lead. lep.}} > 300$  GeV. For the SM-EFT interpretation, the predicted cross-section in each bin is parameterized in terms of Wilson coefficients.

### 13.1.3 Orthogonality considerations

The signal regions of the two measurements use orthogonal event selections:  $m_{e\mu} > 55$  GeV is required in the  $WW$  measurement and  $m_{e\mu} < 55$  GeV in the  $H \rightarrow WW^*$  measurement. However, the latter uses events with  $m_{e\mu} > 55$  GeV to estimate the background from on-shell  $WW$  production, split into two control regions for events with zero and one reconstructed hadronic jet, respectively. The  $WW$  measurement overlaps with the 0-jet control region and, despite the jet-veto requirement, with the 1-jet control region since the  $H \rightarrow WW^*$  and  $WW$  measurements use different jet object definitions:  $p_T^{\text{jet}} > 30$  GeV ( $H \rightarrow WW^*$ ) and  $p_T^{\text{jet}} > 35$  GeV ( $WW$ ), respectively.



**Figure 13.2:** The measured fiducial cross-section of inclusive  $WW \rightarrow e^\pm \mu^\mp$  production for the  $p_T^{\text{lead, lep}}$  (left) and  $m_{e\mu}$  (right) distributions after unfolding is compared with the NNLO prediction (including NLO EW and NLO  $gg \rightarrow WW$  corrections), and NLO + parton shower predictions from Powheg-Box+Pythia8, Powheg-Box+Herwig++, and Sherpa for  $qq$ -initiated  $WW$ , combined with Sherpa+OpenLoops (LO + parton shower) for  $gg$  initial states. All  $qq$  NLO + parton shower predictions are normalized to the NNLO theoretical prediction for the total cross-section, with the  $gg$  LO +parton shower contribution normalized to NLO. Theoretical predictions include PDF and scale uncertainties [259].

For the combination, these control regions of the  $H \rightarrow WW^*$  analysis are removed, and the signal region of the  $WW$  measurement is used instead to constrain the  $WW$  background yield in the  $H \rightarrow WW^*$  signal region; this method avoids double counting of data. The removal of these control regions, with the addition of the  $WW$  measurements, results in only a moderate decrease in sensitivity of the  $H \rightarrow WW^*$  measurement. Therefore, for the combination of the two measurements, the  $WW$  measurement is used to control the uncertainties on the  $WW$  background to  $H \rightarrow WW^*$ . Uncertainties on this background estimation are accounted for by the statistical model introduced in the uncertainty section.

### 13.2 Simulating the impact of dimension-6 SM-EFT operators

The effect of dimension-6 operators on signal production processes, including  $WW$  production, as well as ggF and VBF Higgs boson production, is studied at the lowest order of perturbation theory where the process occurs in the SM. A  $U(3)^5$  flavour symmetry is assumed, and the  $m_W, m_Z, G_F$  input parameter scheme is used, similar to the basis in Chapter 12.

A total of 24 CP-even operators are identified that affect the signal processes of the two measurements, with each operator having a relative impact of at least 1%, for a unit change in coefficient. The operators and their corresponding Wilson coefficients are summarized

in Tab. 13.1. To illustrate the impact of operators on the separate measurements, they are grouped into categories based on the processes they affect: ggF and VBF Higgs boson production,  $WW$  production, decay of the Higgs boson into  $e^\pm \nu_e \mu^\mp \nu_\mu$ , and the total Higgs width  $\Gamma_H^{\text{tot}}$ .

**Table 13.1:** Dimension-6 Warsaw basis operators affecting the signals of interest, assuming  $U(3)^5$  flavour symmetry. The leftmost column shows the name of the Wilson coefficient used to refer to it in the rest of this chapter. The second column shows the operators corresponding to the Wilson coefficient. The right columns indicate which processes or parameters are affected by the operator.

| Wilson coefficient and operator |                                                                             | Affected Processes          |                      |                    |                                       |                         |
|---------------------------------|-----------------------------------------------------------------------------|-----------------------------|----------------------|--------------------|---------------------------------------|-------------------------|
| $c_i^{(6)}$                     | $\mathcal{O}_i^{(6)}$                                                       | $qq \rightarrow e\nu\mu\nu$ | $qq \rightarrow Hqq$ | $gg \rightarrow H$ | $\Gamma_{H \rightarrow e\mu\nu\nu}^f$ | $\Gamma_H^{\text{tot}}$ |
| $c_W$                           | $\epsilon^{IJK} W_\mu^{I\nu} W_\nu^{J\rho} W_\rho^{K\mu}$                   | ✓                           |                      |                    |                                       | ✓                       |
| $c_{HW}$                        | $H^\dagger H W_{\mu\nu}^I W^{I\mu\nu}$                                      |                             | ✓                    |                    | ✓                                     | ✓                       |
| $c_{HWB}$                       | $(H^\dagger \tau^I H W_{\mu\nu}^I B^{\mu\nu})$                              | ✓                           | ✓                    |                    |                                       | ✓                       |
| $c_{HB}$                        | $H^\dagger H B_{\mu\nu} B^{\mu\nu}$                                         |                             | ✓                    |                    |                                       | ✓                       |
| $c_{HDD}$                       | $(H^\dagger D^\mu H)^* (H^\dagger D_\mu H)$                                 | ✓                           | ✓                    | ✓                  | ✓                                     | ✓                       |
| $c_{H\Box}$                     | $(H^\dagger H) \Box (H^\dagger H)$                                          |                             | ✓                    | ✓                  | ✓                                     | ✓                       |
| $c_{H\tilde{G}}$                | $H^\dagger H G_{\mu\nu}^A G^{A\mu\nu}$                                      |                             |                      | ✓                  |                                       | ✓                       |
| $c_{dH}$                        | $(H^\dagger H)(\bar{q}_p d_r H)$                                            |                             |                      |                    |                                       | ✓                       |
| $c_{uH}$                        | $(H^\dagger H)(\bar{q}_p u_r H)$                                            |                             |                      | ✓                  |                                       | ✓                       |
| $c_{eH}$                        | $(H^\dagger H)(\bar{l}_p e_r H)$                                            |                             |                      |                    |                                       | ✓                       |
| $c_{uG}$                        | $(\bar{q}_p \sigma^{\mu\nu} T^A u_r)(\tilde{H} G_{\mu\nu}^A)$               |                             |                      | ✓                  |                                       | ✓                       |
| $c_{uW}$                        | $(\bar{q}_p \sigma^{\mu\nu} u_r)(\tau^I \tilde{H} W_{\mu\nu}^I)$            |                             |                      |                    |                                       | ✓                       |
| $c_{uB}$                        | $(\bar{q}_p \sigma^{\mu\nu} u_r)(\tilde{H} B_{\mu\nu})$                     |                             |                      |                    |                                       | ✓                       |
| $c_{Hd}$                        | $(H^\dagger i \overleftrightarrow{D}_\mu^I H)(\bar{d} \gamma^\mu d)$        | ✓                           | ✓                    |                    |                                       |                         |
| $c_{Hu}$                        | $(H^\dagger i \overleftrightarrow{D}_\mu^I H)(\bar{u} \gamma^\mu u)$        | ✓                           | ✓                    |                    |                                       |                         |
| $c_{Hq}^{(1)}$                  | $(H^\dagger i \overleftrightarrow{D}_\mu^I H)(\bar{q} \gamma^\mu q)$        | ✓                           | ✓                    |                    |                                       |                         |
| $c_{Hq}^{(3)}$                  | $(H^\dagger i \overleftrightarrow{D}_\mu^I H)(\bar{q} \tau^I \gamma^\mu q)$ | ✓                           | ✓                    |                    |                                       | ✓                       |
| $c_{Hl}^{(1)}$                  | $(H^\dagger i \overleftrightarrow{D}_\mu^I H)(\bar{l} \gamma^\mu l)$        | ✓                           |                      |                    |                                       |                         |
| $c_{Hl}^{(3)}$                  | $(H^\dagger i \overleftrightarrow{D}_\mu^I H)(\bar{l} \tau^I \gamma^\mu l)$ | ✓                           | ✓                    | ✓                  | ✓                                     | ✓                       |
| $c_{ll}^{(1)}$                  | $(\bar{l}_p \gamma_\mu l_r)(\bar{l}_s \gamma^\mu l_t)$                      | ✓                           | ✓                    | ✓                  | ✓                                     | ✓                       |
| $c_{ld}$                        | $(\bar{l}_p \gamma_\mu l_r)(\bar{d}_s \gamma^\mu d_t)$                      | ✓                           |                      |                    |                                       |                         |
| $c_{lq}^{(1)}$                  | $(\bar{l}_p \gamma_\mu l_r)(\bar{q}_s \gamma^\mu q_t)$                      | ✓                           |                      |                    |                                       |                         |
| $c_{lq}^{(3)}$                  | $(\bar{l}_p \gamma_\mu \tau^I l_r)(\bar{q}_s \gamma^\mu \tau^I q_t)$        | ✓                           |                      |                    |                                       |                         |
| $c_{lu}$                        | $(\bar{l}_p \gamma_\mu l_r)(\bar{u}_s \gamma^\mu u_t)$                      | ✓                           |                      |                    |                                       |                         |

Similar to the samples generated in § 12.4.1, the impact of each dimension-6 operator's interference with the SM is generated using MadGraph5\_aMC@NLO [109]. The effects of dimension-6 operators contributing to  $q\bar{q} \rightarrow e^\pm \nu_e \mu^\mp \nu_\mu$  production are simulated at leading-order QCD using the SMEFTsimv2.1 [260] model and MadGraph5\_aMC@NLO version 2.8.1 [109], interfaced with PYTHIA 8 [110] using the A14 tune [113] for parton showering and hadronization.

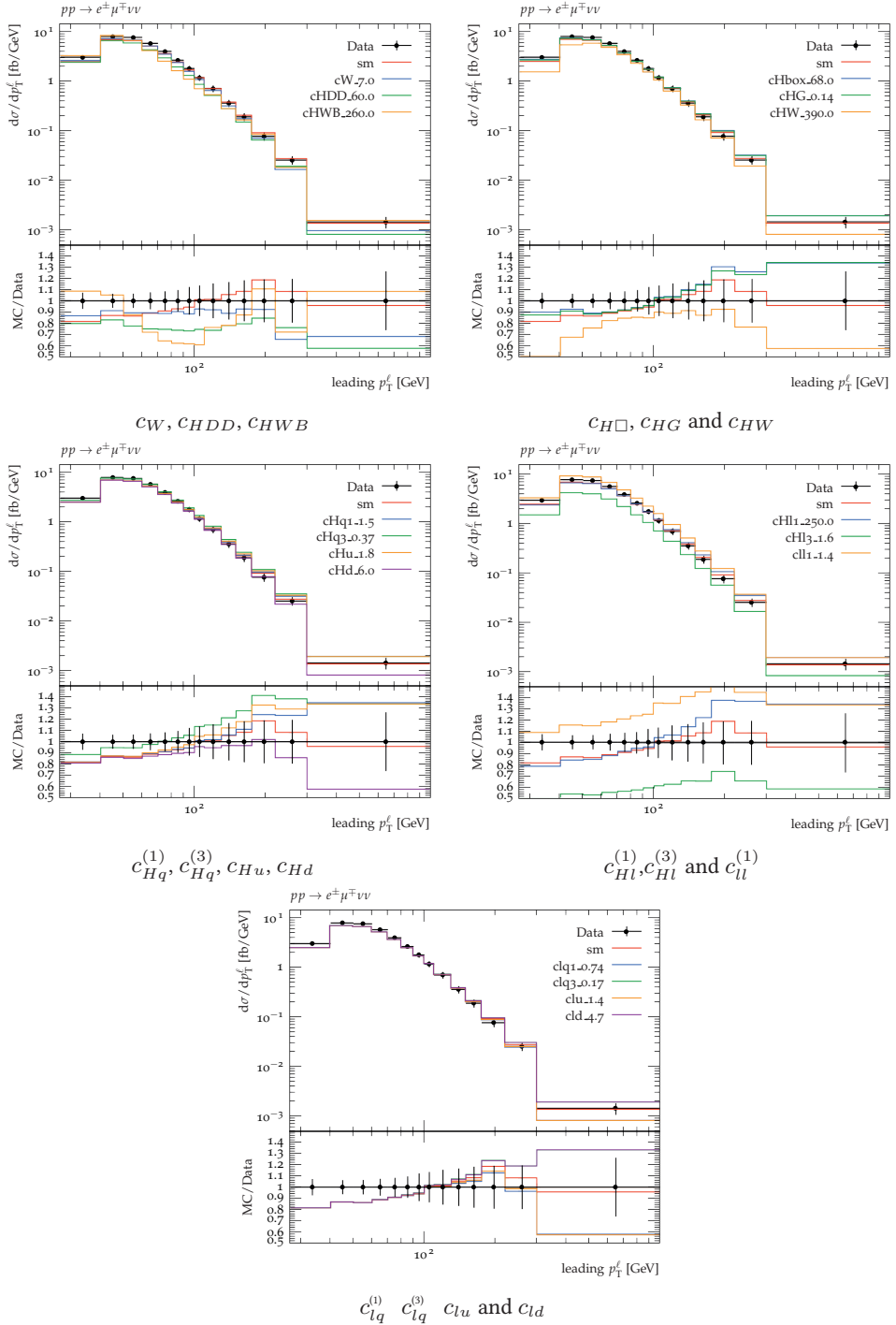
Zero and one-jet events are simulated separately in MadGraph5\_aMC@NLO, and the overlap between matrix element and parton shower emissions is removed using the CKKW-L merging procedure [261, 262]. In the signal region of the  $WW$  analysis, the impact of operators is parameterized differentially across  $p_T^{\text{lead. lep.}}$  analysis bins.

In the phase-space of the  $WW$  analysis, the loop-induced  $gg \rightarrow \mu^\pm e^\mp \nu \nu$  process is simulated using SMEFT@NLO for all operators affecting ggF production and Higgs boson decays into  $e^\pm \nu_e \mu^\mp \nu_\mu$ . While this process constitutes a small correction to  $WW$  production, its consistent treatment alongside on-shell and off-shell Higgs boson production is crucial for accurately modelling operator degeneracies involved in ggF production. Operators not affecting Higgs boson production are excluded from consideration for  $gg \rightarrow \mu^\pm e^\mp \nu \nu$ , as they primarily represent a higher-order QCD correction to  $WW$  production.

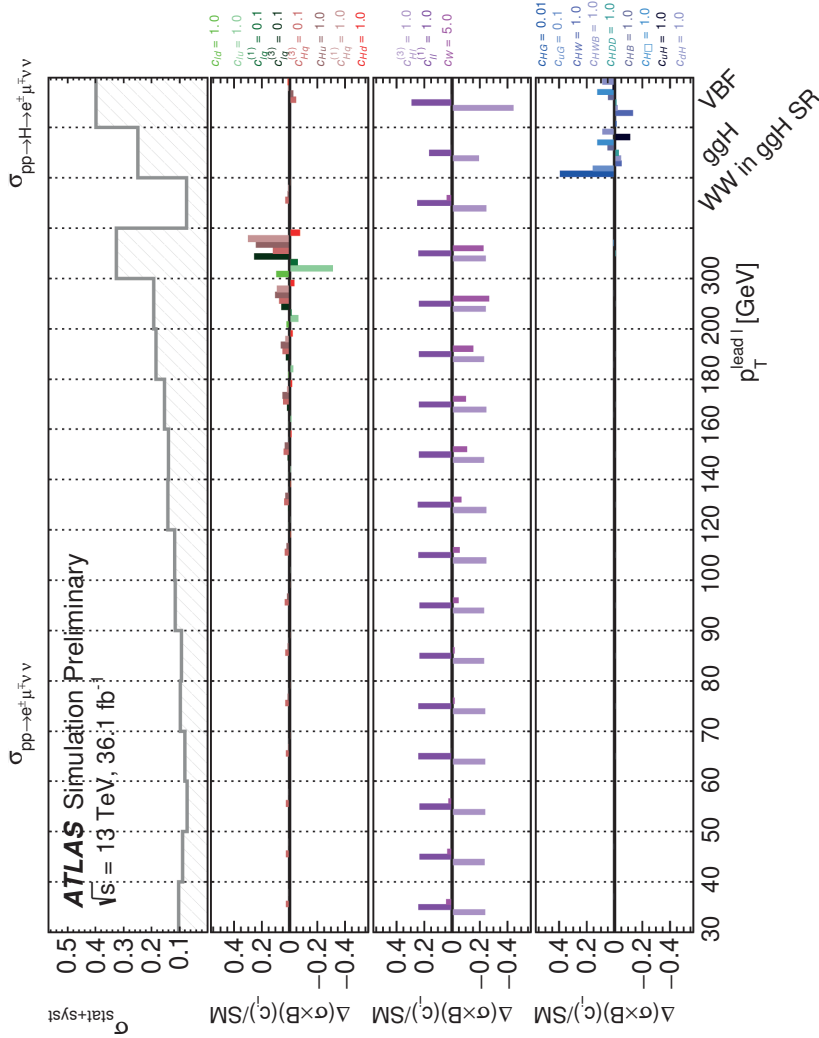
In the signal region of the  $H \rightarrow WW^*$  analysis, ggF Higgs boson production is simulated separately from the  $gg \rightarrow \mu^\pm e^\mp \nu \nu$  background, since the interference among the two processes is negligible in the on-shell regions. Given  $gg \rightarrow \mu^\pm e^\mp \nu \nu$  constitutes only 10% of the total  $WW$  background and can be treated as a higher-order correction to  $WW$  production, SM-EFT effects on this process are neglected in the signal region of the  $H \rightarrow WW^*$  measurement.

The predicted differential cross-sections for the  $WW$  analysis are rescaled in each bin by the ratio of the highest-order particle-level prediction available to that obtained from MadGraph5\_aMC@NLO+ PYTHIA 8. In this rescaling,  $q\bar{q} \rightarrow e^\pm \nu_e \mu^\mp \nu_\mu$  and  $gg \rightarrow e^\pm \nu_e \mu^\mp \nu_\mu$  are treated separately before being combined in the statistical framework. Fig. 13.3 shows the simulation of the SM along with the most relevant  $\mathcal{O}_i^{(6)}$  impact in the  $WW$  analysis, compared to data and the SM distribution. The SM-EFT parameterization of the differential  $pp \rightarrow e^\pm \nu_e \mu^\mp \nu_\mu$  cross-section in the  $WW$  analysis and the ggF and VBF signal strength times branching ratio is summarized in Fig. 13.4. The impact of SM-EFT operators in the background normalization in the  $H \rightarrow WW^* WW$  control region is also shown and consistently taken into account in both measurements.





**Figure 13.3:** The impact of various SM-EFT operators on the  $p_T^{\text{lead, lep.}}$  distribution turned on one-at-a-time, generated with MadGraph5\_aMC@NLO+ PYTHIA 8. The distributions are normalized to the SM NNLO theoretical prediction using MATRIX [94, 95]. They are compared to data, available at [263], and the (NNLO) SM prediction.



**Figure 13.4:** The linear impact of the SM-EFT operators on the signal strength modifiers of the  $H \rightarrow WW^*$  analysis, the normalization of the  $WW$  background in the ggF signal region, and the differential cross-section measured in the  $WW$  analysis, relative to the SM cross-section. In order to illustrate experimental sensitivity to these effects in data, the expected total uncertainty on each modifier is shown in the top panel.  $\Lambda = 1$  TeV is assumed and the values of Wilson coefficients are chosen such that the impact of all operators is of similar size, for better visibility.

### 13.3 The statistical model

Combining measurements from different analysis techniques requires additional effort to ensure consistency. This section provides a brief recap of the statistical models used for both measurements and explains their combined approach.

#### 13.3.1 The statistical model of the $H \rightarrow WW^*$ measurement

The likelihood for interpreting the  $H \rightarrow WW^*$  signal strength measurement follows a similar definition as used in § 12.4.3. However, in this case, the likelihood uses a granularity where all STXS regions are combined into one inclusive STXS region per production mode. In this case, the likelihood is given by

$$L^{HWW}(N|\mathbf{c}, \boldsymbol{\theta}^{HWW}) = \prod_b^{n_{\text{bins}}^{\text{SR+CR}}} \text{Poisson} \left( N_b \left| \sum_p \mu_p(\mathbf{c}) S_{p,b}(\boldsymbol{\theta}^{HWW}) + B_b(\mathbf{c}, \boldsymbol{\theta}^{HWW}) \right. \right) \\ \times \prod_i^{n_{\text{syst}}^{HWW}} f_i(\theta_i^{HWW}),$$

where the background  $B_b$  in this context depends not only on nuisance parameters but also on Wilson coefficients, as listed in Tab. 13.1; their impact on the  $WW$  background across all regions of the  $H \rightarrow WW^*$  analysis is accounted for. In this analysis, the  $WW$  background is determined in the same fit by the incorporation of a control region predominantly sensitive to  $WW$  production. Specifically, the ggF  $N_{\text{jet}} = 0$  and ggF  $N_{\text{jet}} = 1$  regions feature separate strength parameters labelled  $WW$  0-jet and  $WW$  1-jet, scaling this process.

#### 13.3.2 The statistical model of the $WW$ measurement

Because of the nature of the  $WW$  analysis and its fiducial cross-section measurement, the probability density function in the interpretation of the unfolded  $WW$  measurement approach, is based on a multivariate Gaussian distribution for the interpretation. Experimental and theoretical uncertainties are modelled using nuisance parameters, similar to the  $H \rightarrow WW^*$  analysis, and are constrained by Gaussian distributions that represent auxiliary measurements as well as theoretical expectations. The prediction of the analysed  $p_{\text{T}}^{\text{lead, lep.}}$  distribution depends on the set of Wilson coefficients  $\mathbf{c}$ , listed in Tab. 13.1, and is subject to theory systematics. For a given theory nuisance parameter configuration  $\boldsymbol{\theta}_{\text{theo syst}}^{WW} = (\theta_1, \dots, \theta_j, \dots, \theta_{n_{\text{theo syst}}^{WW}})$ , where  $n_{\text{theo syst}}^{WW}$  is the number of theory nuisance parameters affecting the signal predictions, the predicted fiducial cross-section in a bin  $b$  of the unfolded  $p_{\text{T}}^{\text{lead, lep.}}$  distribution is parameterized as:

$$x_b^{\text{pred}}(\mathbf{c}, \boldsymbol{\theta}_{\text{theo syst}}^{WW}) = x_b^{\text{SM}} \left( 1 + \sum_i c_i A_{b,i} \right) \prod_j^{n_{\text{theo syst}}^{WW}} (1 + \theta_j u_{b,j}),$$

where  $u_{b,j}$  represents the relative size of uncertainty  $j$  in bin  $b$ . This formulation assumes the nature and relative effect of theory uncertainties is consistent for both linear EFT effects and the SM.

The measured cross-section depends on the model used for background subtraction, calculation of detector efficiencies, and distribution unfolding. Uncertainties in this model are parameterized using experimental nuisance parameters. Similar to the theoretical uncertainties, they are defined as  $\theta_{\text{exp syst}}^{WW} = (\theta_1, \dots, \theta_j, \dots, \theta_{n_{\text{exp syst}}^{WW}})$ , where  $n_{\text{exp syst}}^{WW}$  is the number of experimental nuisance parameters affecting the measurement. This information is used to cross-section variations corresponding to  $1\sigma$  variations of the nuisance parameter values:

$$x_b^{\text{measurement}}(\theta_{\text{exp syst}}^{WW}) = x_b \times \prod_j^{n_{\text{exp syst}}^{WW}} (1 + \theta_j v_{b,j}),$$

where  $x_b$  is the nominal measurement result and  $v_{b,j}$  is the relative impact of the systematic variation  $j$  on the measured cross-section in bin  $b$ . The nominal measured differential cross-section, in  $n_{\text{bins}}^{WW} = 14$  bins, is given by  $\mathbf{x}$ . Statistical uncertainties of the unfolded data are correlated and encoded in a covariance matrix  $C$ .

The likelihood  $L^{WW}(\mathbf{x}|\mathbf{c}, \theta^{WW})$ , which is modelled as a multivariate Gaussian distribution and is a function of the Wilson coefficients  $\mathbf{c}$ , as well as the union of experimental and theory nuisance parameters,  $\theta^{WW}$ , is then given by

$$L^{WW}(\mathbf{x}|\mathbf{c}, \theta^{WW}) = \frac{1}{\sqrt{(2\pi)^{n_{\text{bins}}^{WW}} \det(C)}} \exp\left(-\frac{1}{2} \Delta \mathbf{x}^\top(\mathbf{c}, \theta^{WW}) C^{-1} \Delta \mathbf{x}(\mathbf{c}, \theta^{WW})\right) \times \prod_i^{n_{\text{syst}}^{WW}} f_i(\theta_i^{WW}).$$

Here  $n_{\text{syst}}^{WW}$  is the total number of experimental and theory nuisance parameters,  $f_i$  correspond to the Gaussian constraints on nuisance parameters, the vector  $\Delta \mathbf{x}$  represents the difference between measurement and prediction, its components  $\Delta x_b$  are defined as

$$\Delta x_b(\mathbf{c}, \theta^{WW}) = x_b^{\text{measurement}}(\theta_{\text{exp syst}}^{WW}) - x_b^{\text{pred}}(\mathbf{c}, \theta_{\text{theo syst}}^{WW}).$$

### 13.3.3 The combined statistical model

The likelihoods from the  $WW$  and  $H \rightarrow WW^*$  measurements are combined to perform a coherent EFT analysis of both measurements. Systematic uncertainties due to identical sources are parameterized with the same nuisance parameter in both measurements, so that the total number of nuisance parameters  $n_{\text{syst}}^{\text{comb}}$  is less than the sum of  $n_{\text{syst}}^{WW}$  and  $n_{\text{syst}}^{HWW}$ . In total 332, nuisance parameters affect the interpretation of both measurements, while 285 affect only the  $H \rightarrow WW^*$  and 137 only the  $WW$  measurement. The  $WW$

background normalization is modelled analogously to the signal parameterization in the  $WW$  measurement, which is a product of the SM prediction, its uncertainty model, and an EFT parameterization. Uncertainties affecting  $WW$  production are treated as fully correlated between the  $WW$  and  $H \rightarrow WW^*$  measurements, such that inclusion of the  $WW$  measurement in the fit reduces the uncertainty on the  $WW$  background yield in the  $H \rightarrow WW^*$  signal region.

The combined likelihood, for an observation of  $N$  events in the  $H \rightarrow WW^*$  analysis bins and a measured differential cross-section of  $\mathbf{x}$  in the  $WW$  analysis, as a function of the vector of Wilson coefficients  $\mathbf{c}$ , is then given by

$$\begin{aligned}
 L^{\text{comb}}(N, \mathbf{x} | \mathbf{c}, \boldsymbol{\theta}^{\text{comb}}) &= \prod_{b \notin WW\text{-CR}}^{n_{\text{bins}}^{\text{SR+CR}}} \text{Poisson} \left( N_b \left| \sum_p \mu_p(\mathbf{c}) S_{p,b}(\boldsymbol{\theta}^{HWW}) + B_b(\mathbf{c}, \boldsymbol{\theta}^{HWW}) \right. \right) \\
 &\times \frac{1}{\sqrt{(2\pi)^{n_{\text{bins}}^{WW}} \det(C)}} \\
 &\times \exp \left( -\frac{1}{2} \Delta \mathbf{x}(\mathbf{c}, \boldsymbol{\theta}^{WW})^\top C^{-1} \Delta \mathbf{x}(\mathbf{c}, \boldsymbol{\theta}^{WW}) \right) \\
 &\times \prod_i^{n_{\text{syst}}^{\text{comb}}} f_i(\theta_i^{\text{comb}}).
 \end{aligned}$$

The best-fit values of the Wilson coefficients are determined by maximizing the combined likelihood. The confidence intervals are estimated using Wilks' theorem as applied in § 12.4.3.

### 13.4 Systematics

Both measurements are subject to similar systematic uncertainties which arise from the reconstruction of jets,  $W$  bosons,  $E_T^{\text{miss}}$  and suppression of the top quark background. The dominant experimental uncertainty sources arise from the modelling of the JES and JER,  $b$ -tagging efficiency, luminosity measurement, pile-up, as well as lepton identification and isolation efficiency. The main background uncertainties are related to the modelling of the  $t\bar{t}$  and single-top quark background, as well as the background due to non-prompt and mis-identified leptons. All dominant uncertainty sources listed above are treated as correlated between measurements, except for those related to the non-prompt background estimate, as different estimation methods are used.

Theoretical uncertainties in the modelling of signal processes also play an important role as the two analyses target a different but still similar topology, therefore the modelling has to be performed coherently. They are evaluated for the SM processes, where it is assumed that the relative uncertainty in each analysis bin remains valid in the presence of dimension-6 EFT operators. The uncertainty of the inclusive  $q\bar{q} \rightarrow e^\pm \nu_e \mu^\mp \nu_\mu$  production cross-section, which is calculated at NNLO QCD with the MATRIX [94, 95] framework, is estimated by varying factorization and renormalization scales by a factor of two, by

taking into account the NNPDF3.0NNLO uncertainty, and by comparing the MMHT2014 [214] and CT14nnlo [264] PDF sets. The inclusive uncertainty is found to be 4.3% and is subsequently applied to the  $WW$  measurement. This inclusion of this uncertainty is necessary as this measurement previously measured the cross-section and now makes use of this information in the reinterpretation.

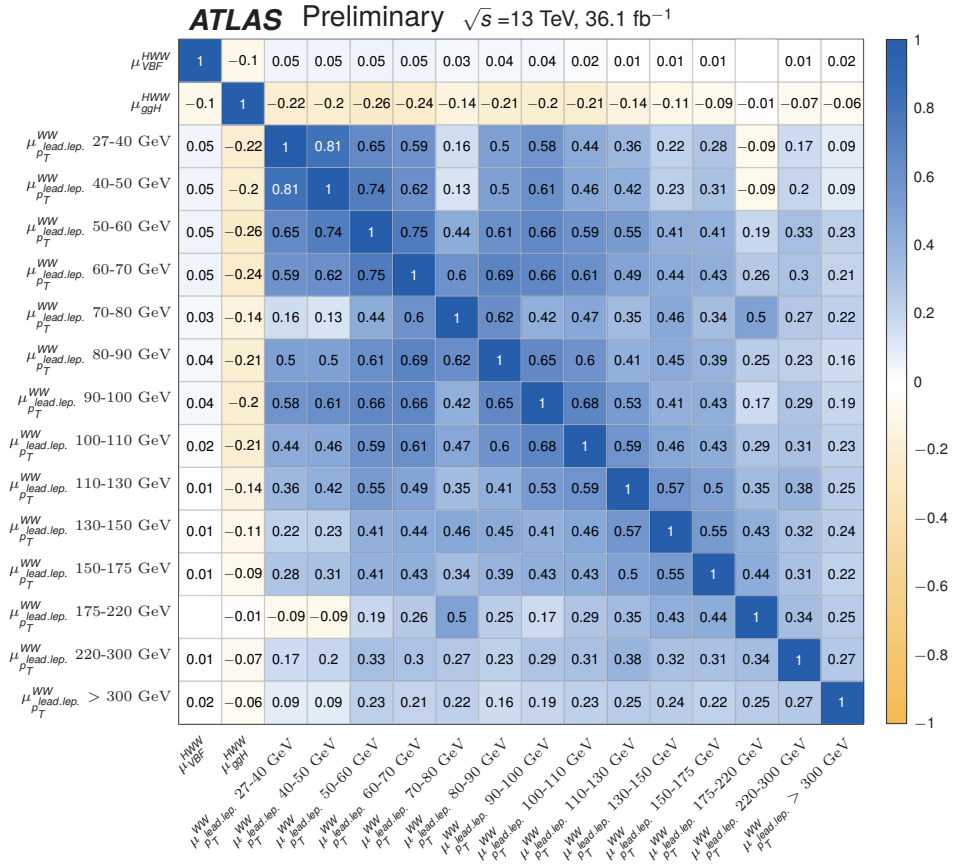
Additionally, shape and acceptance uncertainties on the particle-level prediction of  $q\bar{q} \rightarrow e^\pm \nu_e \mu^\mp \nu_\mu$  production are taken into account. These arise due to the choice of factorization and normalization scale, the PDF, matrix element generator, parton shower, and the matrix element parton shower merging. Uncertainties are evaluated coherently for predictions in the  $WW$  and  $H \rightarrow WW^*$  analysis phase-space.

Modelling uncertainties are treated as correlated between the  $H \rightarrow WW^*$  and  $WW$  measurements. This allows to take into account the theoretical uncertainties when extrapolating the  $pp \rightarrow WW$  process from the  $WW$  analysis to the  $H \rightarrow WW^*$  signal regions. For Higgs boson production, off-shell Higgs,  $gg \rightarrow WW$  production, and other background processes, the same uncertainties as already evaluated for the  $WW$  and  $H \rightarrow WW^*$  measurements are used.

### 13.5 Signal strength and Wilson coefficient measurements

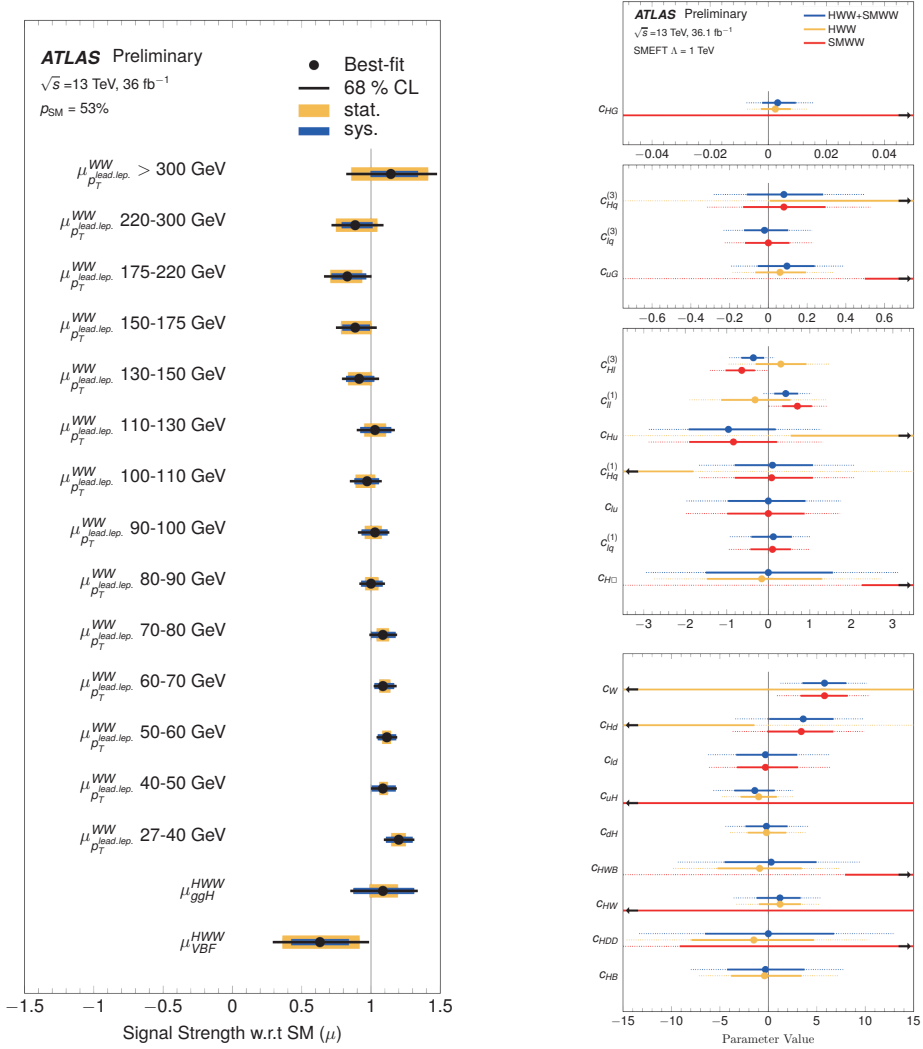
In order to determine the deviations between SM predictions and measurements in each region, a simultaneous measurement of 16 signal strength parameters is performed. Analogously to the STXS SM-EFT interpretation in Chapter 12, this can be viewed as the input to the interpretation. These inputs are the two signal strength parameters of the  $H \rightarrow WW^*$  analysis and 14 signal strength parameters which scale the predicted SM cross-section in each  $p_T^{\text{lead. lep.}}$  bin of the  $WW$  analysis. This measurement is first performed using a likelihood function that is identical to the combined likelihood  $L^{\text{comb}}(\mathbf{N}, \mathbf{x} | \mathbf{c}, \boldsymbol{\theta}^{\text{comb}})$ , except that each per-bin signal strength is parameterized with a single dedicated parameter.

The measurements' observed correlations and signal strength results are shown in Figs. 13.5 and 13.6, respectively. The signal strength results are presented, along with the the 68% and 95% confidence intervals limits on each Wilson coefficient  $c_i$ . The combined measurements' central values differ by less than 2% from the stand-alone measurement. After combining, a degradation of 10% in total uncertainty for the ggF signal strength measurement is introduced due to the replacement of the  $WW$  control region with the  $WW$  measurement, whereas the VBF measurement precision is unaffected. This is due to the  $WW$  differential measurement being slightly worse in determining the overall normalization of the  $WW$  prediction because of the differential nature of the measurement.



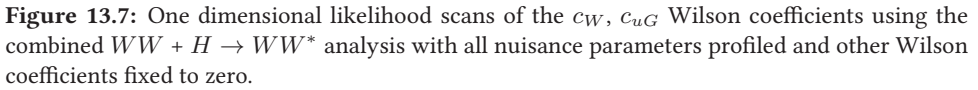
**Figure 13.5:** Observed correlation matrix of the signal strength modifiers of the  $H \rightarrow WW^*$  analysis and the  $WW$  measurement.

The  $WW$  signal strengths, in bins of  $p_T^{\text{lead. lep.}}$ , are similar to the ratio of unfolded data to the SM prediction that is presented in the original measurement. Differences in this ratio can arise because nuisance parameters which parameterize uncertainties on the measurement, or on the differential  $WW$  cross-section prediction, are pulled from their nominal value in the maximum likelihood fit. A selection of the likelihood scans performed on the full combination, from which the projected Wilson coefficient limits are extracted, are shown in Fig. 13.7.

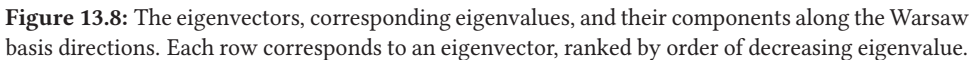


**Figure 13.6:** Measured signal strengths from the combined fit of signal-strength modifiers in bins of the  $WW$  analysis and for the two signal strength modifiers of the  $H \rightarrow WW^*$  analysis. Statistical uncertainties are shown in yellow, systematic uncertainties as blue bars and the total uncertainties are indicated by black lines. The compatibility of the observed data with the SM hypothesis ( $p_{\text{SM}}$ ) is 53%. (left) Observed limits on individual Wilson coefficients, while setting other Wilson coefficients to zero. Limits observed from an analysis of the SM  $WW$ , the  $H \rightarrow WW^*$ , and the combined measurements are compared. The stand-alone  $H \rightarrow WW^*$  analysis is not identical to the one that enters the combination, where the  $WW$  0-jet and 1-jet control regions are removed. 68% (95%) confidence level intervals are shown as solid (dotted) lines. Portions of the confidence intervals outside the boundaries of the figure are clipped and different parameter value axis limits are used for illustrative purposes. If a measurement has no sensitivity to a Wilson coefficient, no interval is drawn.  $\Lambda = 1 \text{ TeV}$  is assumed (right).





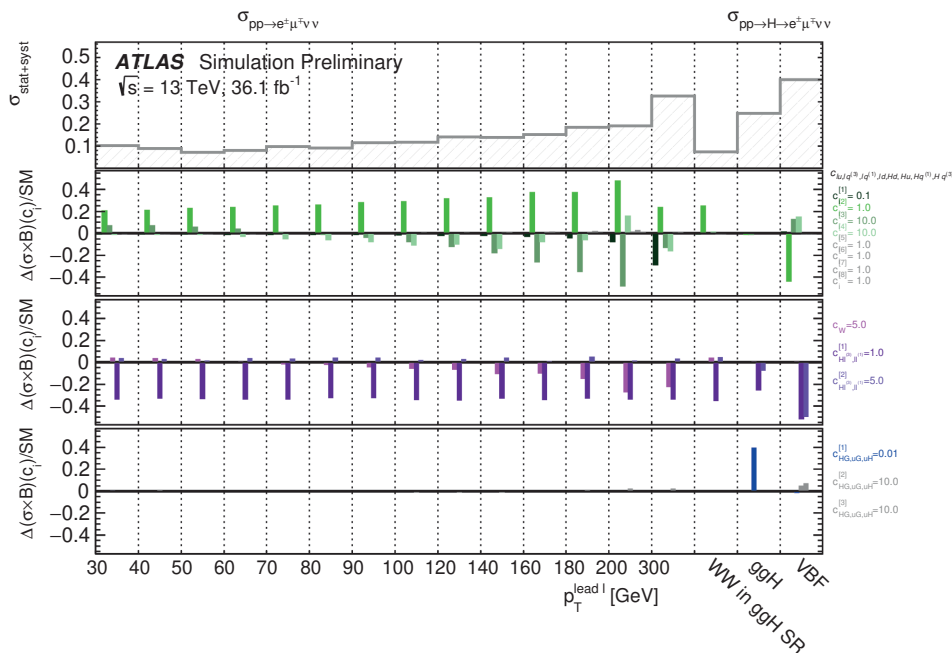
Similar to the analysis of [Chapter 12](#), flat directions in the parameter space must be identified, and to construct the fit model, an orthogonal basis of directions is defined among the remaining parameters. [Fig. 13.8](#) lists eigenvectors  $c_j$  and the eigenvalues of the Fisher information matrix. The first eight directions are shown, which have an eigenvalue of  $\lambda \gtrsim 0.01$ . These eigenvectors, each probing linear combinations of a large numbers of Wilson coefficients, are not directly measured. Instead, they provide a measure of how many sensitive directions the combination of analyses can constrain.



In order to improve the future interpretability of the measurement, parameter groups containing Wilson coefficients with a similar physics impact are formed. The operator  $c_W$  is treated as its own subgroup, as its physics effect uniquely only affects weak boson self-couplings. The operators  $c_{HG}$ ,  $c_{uG}$ , and  $c_{uH}$  form a group, as these operators affect

ggF Higgs production as a function of  $p_T^H$ . The set of  $c_{Hl}^{(3)}$  and  $c_u^{(1)}$  form a group, as both lead to a shift of the Fermi constant  $G_F$  and mainly change the normalization of the observed processes.

The coefficients  $c_{lu}$ ,  $c_{lq}^{(3)}$ ,  $c_{lq}^{(1)}$ ,  $c_{ld}$ ,  $c_{Hd}$ ,  $c_{Hu}$ ,  $c_{Hq}^{(1)}$ , and  $c_{Hq}^{(3)}$  are grouped because their most important effect is the modification of  $q\bar{q} \rightarrow e^\pm \nu_e \mu^\mp \nu_\mu$  production at high momentum transfer  $p_T^{\text{lead, lep.}}$ . These operators affect both effective boson-fermion and four-fermion couplings. The next two groups contain  $c_{HB}$ ,  $c_{HW}$  and  $c_{HWB}$  and affect Higgs boson decays into  $e^\pm \nu_e \mu^\mp \nu_\mu$  and VBF production. The impact of the rotated fit parameters  $c_j$  on the measured signal strength is shown in [Fig. 13.9](#).

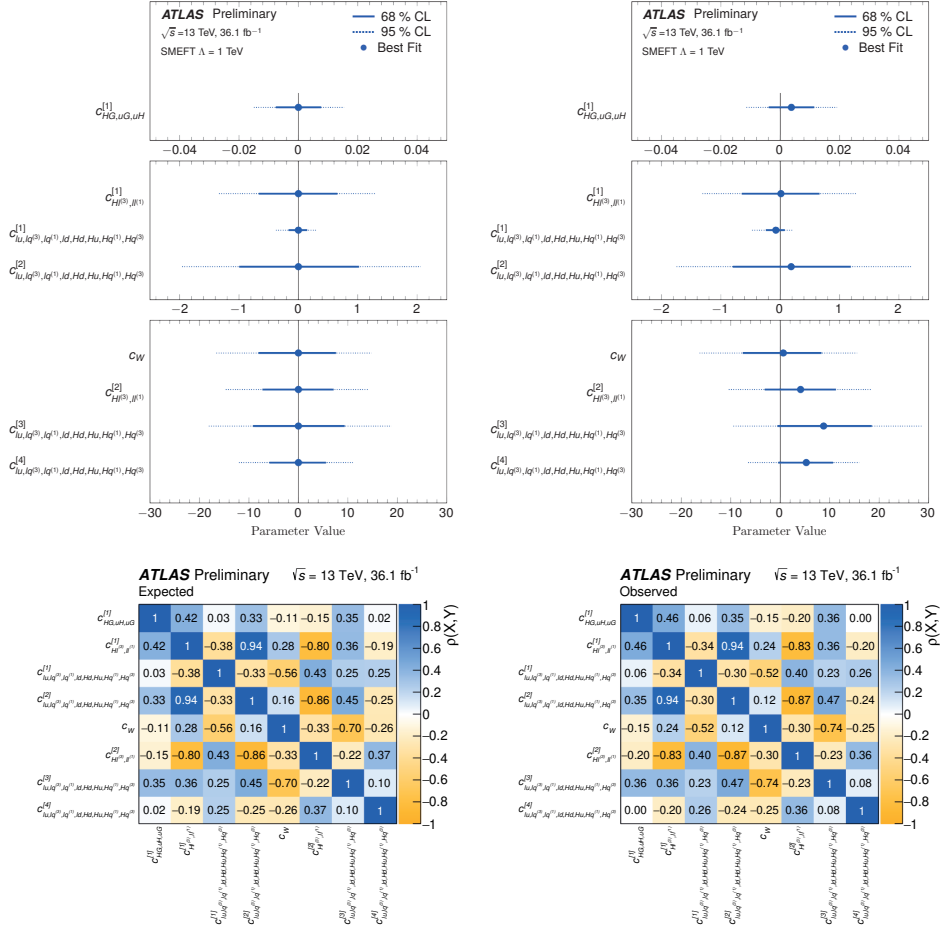


**Figure 13.9:** Impact of the rotated  $c_j$  basis for  $\Lambda = 1$  TeV on the signal strength modifiers of the  $H \rightarrow WW^*$  analysis, the normalization of the  $WW$  background in the ggF signal region, and the differential cross-section measured in the  $WW$  analysis, relative to the SM cross-section. Operators of the groups that are not included in the fit are shown in grey. In order to illustrate experimental sensitivity to constrain the operators from the data, the total expected uncertainty on the corresponding regions  $\sigma_{\text{stat+sys}}$  is shown in the top panel.

### 13.7 SM-EFT interpretation results

In total eight parameters are measured, comprised of seven sensitive directions and Wilson coefficient  $c_W$ . The constraints and correlations of the measured sensitive directions are shown in Fig. 13.10. The first  $c_{H^{(3)}, ll^{(1)}}^{[1]}$  eigenvector changes the normalization of  $WW$  production and affects the VBF and ggF production rates. The second eigenvector in this group,  $c_{H^{(3)}, ll^{(1)}}^{[2]}$ , mainly affects the rate of VBF Higgs boson production, while

$c_{uG,HG,uH}^{[1]}$  modifies the ggF-production rate.

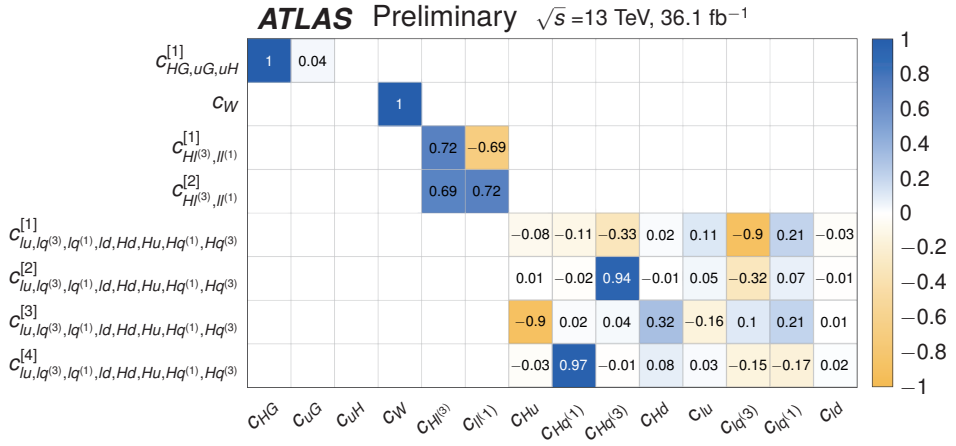


**Figure 13.10:** Summary of expected (left) and observed (right) measurements in the space of the sensitive directions. In the top figures, ranges shown correspond to 68% (solid) and 95% (dashed) confidence level intervals, where all other sensitive directions and nuisance parameters are profiled. Observed correlations between the parameters are shown on the the bottom.

The main effect of  $c_{lu,lq^{(3)},lq^{(1)},ld,Hd,Hu,Hq^{(1)},Hq^{(3)}}^{[1]}$  is a change in the rate of  $WW$  and VBF Higgs boson production, while the remaining operators of that group modify the tail of the  $p_T^{\text{lead, lep.}}$  distribution. No eigenvector of the last two groups in Tab. 13.2, which modify exclusively Higgs production and decay rates, has an expected sensitivity above the threshold of 10, when included in the analysis. It was found that only two degrees of freedom can be constrained exclusively by the  $H \rightarrow WW^*$  measurement, these are  $c_{H^{(3)},ll^{(1)}}^{[2]}$  and  $c_{uG,HG,uH}^{[1]}$ . A summary of the sensitive directions is visualized in Fig. 13.11.

**Table 13.2:** The eigenvectors of the operator subspaces, ordered by eigenvalue. Each eigenvector corresponds to a linear combination of Wilson coefficients. Operator subspaces are separated by horizontal lines in the table. The final column indicates the directions included in the combined fit.

| Parameter                                          | Definition                                                                                                                               | Floated<br>POI |
|----------------------------------------------------|------------------------------------------------------------------------------------------------------------------------------------------|----------------|
| $c_W$                                              | $c_W$                                                                                                                                    | ✓              |
| $c_{HG,uG,uH}^{[1]}$                               | $0.999c_{HG} + 0.039c_{uG} - 0.003c_{uH}$                                                                                                | ✓              |
| $c_{HG,uG,uH}^{[2]}$                               | $0.039c_{HG} - 0.999c_{uG} - 0.126c_{uH}$                                                                                                |                |
| $c_{HG,uG,uH}^{[3]}$                               | $0.008c_{HG} - 0.125c_{uG} + 0.992c_{uH}$                                                                                                |                |
| $c_{lu,lq^{(3)},ld,Hd,Hu,Hq^{(1)},Hq^{(3)}}^{[i]}$ |                                                                                                                                          |                |
| 1                                                  | $-0.903c_{lq}^{(3)} + 0.113c_{lu} - 0.080c_{Hu} + 0.209c_{lq}^{(1)} - 0.033c_{ld} + 0.023c_{Hd} - 0.111c_{Hq}^{(1)} - 0.329c_{Hq}^{(3)}$ | ✓              |
| 2                                                  | $-0.319c_{lq}^{(3)} + 0.045c_{lu} + 0.007c_{Hu} + 0.070c_{lq}^{(1)} - 0.012c_{ld} - 0.006c_{Hd} - 0.023c_{Hq}^{(1)} + 0.944c_{Hq}^{(3)}$ | ✓              |
| 3                                                  | $+0.102c_{lq}^{(3)} + 0.159c_{lu} - 0.903c_{Hu} + 0.206c_{lq}^{(1)} + 0.007c_{ld} - 0.323c_{Hd} - 0.000c_{Hq}^{(1)} + 0.036c_{Hq}^{(3)}$ | ✓              |
| 4                                                  | $-0.147c_{lq}^{(3)} + 0.026c_{lu} - 0.033c_{Hu} - 0.174c_{lq}^{(1)} + 0.015c_{ld} + 0.079c_{Hd} + 0.969c_{Hq}^{(1)} - 0.014c_{Hq}^{(3)}$ | ✓              |
| 5                                                  | $-0.178c_{lq}^{(3)} - 0.029c_{lu} - 0.166c_{Hu} - 0.929c_{lq}^{(1)} + 0.049c_{ld} + 0.168c_{Hd} - 0.213c_{Hq}^{(1)} + 0.008c_{Hq}^{(3)}$ |                |
| 6                                                  | $0.120c_{lq}^{(3)} + 0.897c_{lu} - 0.261c_{Hu} - 0.064c_{lq}^{(1)} - 0.175c_{ld} - 0.279c_{Hd} - 0.001c_{Hq}^{(1)} + 0.000c_{Hq}^{(3)}$  |                |
| 7                                                  | $0.068c_{lq}^{(3)} + 0.333c_{lu} + 0.285c_{Hu} + 0.089c_{lq}^{(1)} - 0.110c_{ld} + 0.884c_{Hd} - 0.043c_{Hq}^{(1)} + 0.002c_{Hq}^{(3)}$  |                |
| 8                                                  | $0.005c_{lq}^{(3)} + 0.207c_{lu} - 0.001c_{Hu} + 0.054c_{lq}^{(1)} + 0.976c_{ld} + 0.038c_{Hd} - 0.014c_{Hq}^{(1)} + 0.001c_{Hq}^{(3)}$  |                |
| $c_{Hl^{(3)},ll^{(1)}}^{[1]}$                      | $-0.695c_{ll}^{(1)} + 0.719c_{Hl}^{(3)}$                                                                                                 | ✓              |
| $c_{Hl^{(3)},ll^{(1)}}^{[2]}$                      | $+0.719c_{ll}^{(1)} + 0.695c_{Hl}^{(3)}$                                                                                                 | ✓              |
| $c_{HW,HWB,HDD}^{[i]}$                             |                                                                                                                                          |                |
| 1                                                  | $-0.262c_{HDD} - 0.888c_{HW} - 0.379c_{HWB}$                                                                                             |                |
| 2                                                  | $-0.576c_{HDD} + 0.458c_{HW} - 0.677c_{HWB}$                                                                                             |                |
| 3                                                  | $-0.775c_{HDD} - 0.041c_{HW} + 0.631c_{HWB}$                                                                                             |                |
| $c_{dH,HB,H\Box}^{[i]}$                            |                                                                                                                                          |                |
| 1                                                  | $-0.543c_{dH} - 0.297c_{HB} - 0.786c_{H\Box}$                                                                                            |                |
| 2                                                  | $0.694c_{dH} + 0.369c_{HB} - 0.619c_{H\Box}$                                                                                             |                |
| 3                                                  | $0.473c_{dH} - 0.881c_{HB} + 0.006c_{H\Box}$                                                                                             |                |



**Figure 13.11:** Visualization of the eight measured linear combinations of Wilson coefficients. The grid shows the contribution of Warsaw basis vectors to the linear combinations. The component of  $c_{uH}$  in  $C_{HG,uG,uH}^{[1]}$  is small (-0.003) and is not shown for clarity.

## PART V

### FINAL REFLECTIONS



## CONCLUSION AND DISCUSSION

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In this chapter, the results and conclusions of the measurements presented in this thesis are discussed and, where relevant, compared to similar studies.

### **Analysis of off-shell Higgs boson production using $WW \rightarrow \ell\nu\ell\nu$ final states**

In this thesis, the off-shell Higgs boson production rate is measured in the ggF and VBF production modes, using the  $H^* \rightarrow WW \rightarrow e\nu\mu\nu$  decay channel. The observed signal strengths are found to be  $\mu_{\text{off-shell}}^{\text{ggF}} = 0.2^{+1.3}_{-0.2}$  and  $\mu_{\text{off-shell}}^{\text{VBF}} = 0.4^{+3.4}_{-0.4}$ . The measurement of the individual production processes reveal they are consistent with each other within their uncertainty. The combined measurement of these two production modes resulted in a best-fit value of  $\mu_{\text{off-shell}} = 0.3^{+0.9}_{-0.3}$  with an observed (expected) 95% CL upper limit at 3.4 (4.4). The results indicate that the measured signal strengths are consistent with the SM predictions, but with an observed deficit in the off-shell signal. This corresponds to an excess in the signal-plus-background prediction with interference (SBI), which is caused by the apparent lack of signal, including signal interference with the continuum background. The low best-fit value causes the observed limits to be more stringent than the expected limits.

In comparison, the previous off-shell Higgs  $WW$  measurement, published in 2015 [200], had an observed  $\mu_{\text{off-shell}}$  upper limit of 17.2 at 95% CL using only the  $WW$  process. The current limit is the most stringent limit obtained using the  $WW$  process, just like the previous result had the most stringent limit at the time. The previous results were made with an at the time unknown K-factor ratio, which was allowed to vary between 0.5 and 2.0. The result of the 2015 off-shell  $WW/ZZ$  combination was a 95% CL upper limit in the range 5.1-8.6, depending on the assumed K-factor ratio, while no lower limit above a statistical significance of  $1\sigma$  could be established. In comparison with the full 2015 combination, this analysis performs significantly better by setting an upper limit of 3.4 times the SM at 95% CL, while also setting a lower limit of 0.01 times the SM at 68% CL. The improved performance arises from a larger dataset, higher collision energy, advancements in theory predictions, and novel analysis techniques. Consequently, background uncertainties are reduced to  $\delta_B \sim \mathcal{O}(10\%)$ , which makes the  $H \rightarrow WW \rightarrow \ell\nu\ell\nu$  channel viable for off-shell Higgs boson measurements and indirect Higgs width constraints.

A second comparison can be performed with the more recent Run-2 results of the off-shell



$ZZ$  analysis performed by ATLAS and CMS [108, 199]. These analyses set limits on the off-shell signal strength of  $[0.15, 2.4]$  and  $[0.0061, 2.0]$  at the 95% CL, with a best-fit value of 1.1 and 0.62, respectively. When comparing these limits to the off-shell result in this thesis, it is clear that the  $WW$  analysis competes with the off-shell  $ZZ$  measurements from ATLAS and CMS. This demonstrates its potential to further constrain the off-shell Higgs boson production rate in a future combined analysis.

Another point is the comparison of best-fit values: ATLAS finds  $\mu_{\text{off-shell}} = 1.1$ , CMS reports  $\mu_{\text{off-shell}} = 0.62$ , and this analysis yields  $\mu_{\text{off-shell}} = 0.3$ . All are statistically compatible, though the lower best-fit here reflects a reduced off-shell production rate due to a signal deficit and the excess interfering background.

### **Combined on-shell and off-shell $H \rightarrow WW$ measurement of the Higgs total width**

A combination is performed in order to probe the width and lifetime of the Higgs boson using the off-shell and on-shell ggF and VBF production rates in the  $H \rightarrow WW^* \rightarrow e\nu\mu\nu$  decay channel. The observed width is measured to be  $\Gamma_H = 0.9^{+3.4}_{-0.9}$  MeV. In this combination, the limits on the Higgs width of  $[0.0, 4.3]$  at 68% CL and an upper limit of 13.1 MeV at 95% CL are set. The uncertainty of the observed lifetime appears to be dominated by the off-shell measurement, due to the comparatively low total uncertainty of the on-shell measurement. It should also be noted that the on-shell signal strengths are relatively close to their SM value, which in combination with the low observed signal strength of the off-shell measurement, results in a Higgs lifetime which is much larger than expected value of  $\tau_H = 1.6 \times 10^{-22}$  s. The results, however, are still consistent with the expected SM width and lifetime within their uncertainties.

Similar to the off-shell production measurement, this analysis outperforms the  $WW/ZZ$  combined width measurement published in 2015 [200], with an observed upper limit of 22.7 MeV at 95% CL.

Insights into the relative performance of the width and lifetime measurement can be obtained by comparing again with the Run-2 results of the off-shell  $ZZ$  analysis performed by ATLAS and CMS [108, 199]. The CMS and ATLAS results are statistically compatible and comparable to the off-shell  $WW$  result present in this section, consisting of a best-fit value of 3.2 MeV and 4.5 MeV with the limits  $[0.5, 8.5]$  and  $[0.5, 10.5]$  at the 95% CL, respectively. It should be noticed that the  $WW$  lifetime measurement upper limit is only  $\sim 10\%$  weaker when compared to the ATLAS  $ZZ$  result, while for the off-shell result this was  $\sim 25\%$ . This effect is caused by the difference in observed signal strength values in the measurements. The different values of the observed on-shell and off-shell signal strengths, notably the on-shell  $WW$  signal strength that is larger than the SM expectation, results in an even more competitive upper limit on the Higgs boson width.

## Combined SM-EFT interpretation of Higgs boson production and decay

In Chapter 12, an interpretation of the combined Higgs boson measurements results, using differential STXS measurements within the framework of the SM-EFT is presented. After the implementation of the acceptance corrections and the principal component analysis, a total of 13 degrees of freedom are fitted simultaneously. This number is expected to grow, once more analyses of decay and production models are included. The results of the analysis show that the data is consistent with the SM predictions.

It is observed that there are some deviations above  $1\sigma$  in the combined fit result. The first of these corresponds to the  $gg \rightarrow H$  production mode. The  $c_{HG,uG,uH}^{[1]}$  direction is fitted at a value below the SM prediction, which corresponds to a lower Higgs production rate. The second and third deviation are found in the  $c_{dH}$  and  $c_{eH}$  directions, respectively. The  $c_{eH}$  coefficient, with a best-fit value of 1.8, corresponds to the Higgs boson  $H \rightarrow \tau\tau$  decay rate and a slightly lower value than the SM prediction. Finally, the  $c_{dH}$  parameter is measured to be 3.2, due to a slight deficit in  $H \rightarrow bb$ , events and corresponds to a decreased coupling strength of the  $H \rightarrow bb$  decay vertex.

The results of the second round of SM-EFT interpretations of the combined Higgs measurements are consistent with those of the previous iteration, but features higher precision in all measured directions, as well as a reduced correlation between the measurement parameters. This is caused by the increased amount of data, via the inclusion of the  $H \rightarrow \tau\tau$ ,  $H \rightarrow bb$  decay modes and the VBF and  $t\bar{t}H$  production modes. The newly added channels also increased the number of probed directions, such as  $c_{dH}, c_{eH}$  and  $c_{top}^{[1]}$ , and this number is expected to grow even further in the future with the inclusion of new channels.

## Combined SM-EFT interpretation of Higgs and $WW$ measurements

In Chapter 13, the first combined EW and Higgs SM-EFT interpretation was performed using the  $WW \rightarrow e\nu\mu\nu$  decay channel: it proposed and applied the methods needed for the combination of a Higgs boson STXS measurement with a differential EW measurement. It also served as a stepping stone towards the global combination which includes Higgs, EW, and EWPO measurements [7].

In total, eight eigenvectors are measured and all are in agreement with the SM within their uncertainty. However, relatively large deviations from zero are observed for  $c_W$ ,  $c_{ll}^{(1)}$  and  $c_{Hl}^{(3)}$  in the analysis of the  $WW$  measurement. The deviations of  $c_{ll}^{(1)}$  and  $c_{Hl}^{(3)}$  can be explained by the fact that the fiducial cross-section measured in the  $WW$  analysis is larger than the SM expectation, as  $c_{ll}^{(1)}$  and  $c_{Hl}^{(3)}$  only change the normalization of the signal in this phase-space region.

A large value of the  $c_W$  coefficient is preferred by the data as it corresponds to the harder  $p_T$  distribution of the leading-lepton that is observed in the data. This harder  $p_T$  spectrum can also be observed in the signal strength results. It is also found that the individual

Wilson coefficient results are a good measure of how sensitive a particular analysis is to a given Wilson coefficient, even though it assumed an over-simplified statistical model, where the other operators are assumed to be absent.

In conclusion, it is found that a combination with additional measurements of single and multiboson production can help to improve the constraints on the remaining operators and reduce degeneracies among the parameters that affect boson-fermion and four-fermion couplings. Combinations with more measurements of Higgs-boson production and decay, performed in [Chapter 12](#), further reduce the number of flat directions in the parameter space probed. In general, when not needing to remove any data for orthogonality reasons, the combination of measurements improves the constraints on the Wilson coefficients shared between the analyses significantly, while also reducing unwanted correlations.

## TECHNICAL SUMMARY AND OUTLOOK

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This chapter presents a technical summary of the key results of this thesis, followed by an outlook on future improvements in off-shell Higgs production, Higgs lifetime, and SM-EFT measurements.

### Technical summary

The first project presented in this thesis is the development and tuning of new  $E_T^{\text{miss}}$  trigger algorithms for Run-3. In this work, the KFMET and jwoj algorithms are developed and tuned for the ATLAS trigger menu in Run-3. The chapter concludes with a performance estimate of the algorithms and shows their improvement, compared to the previous Run-2 algorithm. The results show that further tuning and improving L1 MET triggers enhances performance in high pile-up conditions. It is also the first simple regression machine learning network running real-time during data acquisition at L1, trained on offline data, and provides the best performance in high pile-up conditions of all tested algorithms.

This thesis also presents the first ever measurement of the off-shell Higgs boson production in the  $WW$  decay channel at a centre-of-mass energy of  $\sqrt{s} = 13$  TeV using the ATLAS detector. The measurement is performed using the full Run-2 dataset, made possible by the LHC, and collected with the ATLAS experiment, corresponding to an integrated luminosity of  $140 \text{ fb}^{-1}$ . The measurement is performed on the different and same flavour decay modes and interpreted in terms of the Higgs total width and off-shell/on-shell coupling ratios via a combination with the on-shell  $H \rightarrow WW^*$  measurement.

In order to perform this measurement, a novel analysis strategy is developed which makes use of multiple background estimation methods and categorisation of the interfering off-shell signal and non-interfering background processes. This involved using several neural networks and a  $m_{WW}$  mass proxy in a 2D fit configuration. Additionally, state-of-the-art higher-order QCD correction factors and Monte Carlo samples have been made in collaboration with the off-shell theory community, in order to estimate the one- and two-loop heavy quark contributions and boost the sensitivity. All the various background contaminations in this analysis have been studied, estimated and controlled, resulting in relatively low systematic background uncertainty.

The main results of the off-shell Higgs measurements and interpretations in the  $WW$  decay channel are listed below and summarized in [Tab. 15.1](#):

- An observed (expected) excess of events for the off-shell  $pp \rightarrow H^* \rightarrow WW$  process with a significance of  $0.7\sigma$  ( $1.2\sigma$ ).
- A total off-shell Higgs production signal strength of  $\mu_{\text{off-shell}} = 0.3^{+0.9}_{-0.3}$  with an upper limit of  $\mu_{\text{off-shell}} < 3.4$ , at 95% CL<sup>1</sup>.
- An average Higgs lifetime of  $\tau_H = 7.0^{+82.1}_{-5.5} \times 10^{-22}$  s, corresponding to a Higgs width of  $\Gamma_H = 0.90^{+3.4}_{-0.9}$  MeV at an observed (expected) significance of  $0.7\sigma$  ( $1.2\sigma$ ).
- An upper limit on the total Higgs width  $\Gamma_H$  of 13.1 MeV, at 95% CL<sup>1</sup>.

**Table 15.1:** Results for off-shell signal strengths and  $\Gamma_H$  under different conditions. The POI  $\mu_{\text{off-shell}}$  is presented with  $\mathcal{R}_{VV,gg}^{\text{off-shell}} = \mu_V^{\text{off-shell}} / \mu_F^{\text{off-shell}} = 1$ , while  $\mu_V^{\text{off-shell}}$  and  $\mu_F^{\text{off-shell}}$  are constrained with the other signal strengths floating unconstrained ( $\omega$ ). The Higgs width  $\Gamma_H$  is extracted with the on-shell signal strengths unconstrained. The expected values  $\mu_{\text{off-shell}}$ ,  $\mu_{\text{VBF}}^{\text{off-shell}}$ ,  $\mu_{\text{ggF}}^{\text{off-shell}}$  are unity, and  $\Gamma_H^{SM} = 4.1$  MeV. Only upper 95% constraints are provided where lower limits are undefined.

| Parameter                                     | Condition                                           | Observed             |                     | Expected             |                     |
|-----------------------------------------------|-----------------------------------------------------|----------------------|---------------------|----------------------|---------------------|
|                                               |                                                     | 68% CL               | 95% CL <sup>1</sup> | 68% CL               | 95% CL <sup>1</sup> |
| $\mu_{\text{off-shell}}^{\text{ggF}}$         | $\mu_{\text{off-shell}}^{\text{VBF}} (\omega)$      | $0.2^{+1.3}_{-0.2}$  | $< 3.5$             | $1.0^{+2.5}_{-1.0}$  | $< 5.1$             |
| $\mu_{\text{off-shell}}^{\text{VBF}}$         | $\mu_{\text{off-shell}}^{\text{ggF}} (\omega)$      | $0.4^{+3.4}_{-0.4}$  | $< 6.3$             | $1.0^{+2.9}_{-1.0}$  | $< 6.2$             |
| $\mu_{\text{off-shell}}$                      | $\mathcal{R}_{\text{gg},VV}^{\text{off-shell}} = 1$ | $0.3^{+0.9}_{-0.3}$  | $< 3.4$             | $1.0^{+2.3}_{-1.0}$  | $< 4.4$             |
| $\mathcal{R}_{\text{gg}}^{\text{off-shell}*}$ | $\Gamma_H = \Gamma_H^{SM}$                          | 0.2                  | –                   | –                    | –                   |
| $\mathcal{R}_{VV}^{\text{off-shell}}$         | $\Gamma_H = \Gamma_H^{SM}$                          | $0.6^{+1.6}_{-0.6}$  | $< 2.8$             | $1.0^{+1.2}_{-1.0}$  | $< 2.7$             |
| $\mu_{\Gamma_H}$                              | $\mathcal{R}_{\text{gg},VV}^{\text{off-shell}} = 1$ | $0.2^{+0.8}_{-0.2}$  | $< 3.2$             | $1.0^{+2.0}_{-0.9}$  | $< 4.2$             |
| $\Gamma_H$ [MeV] <sup>†</sup>                 | $\mathcal{R}_{\text{gg},VV}^{\text{off-shell}} = 1$ | $0.9^{+3.4}_{-0.9}$  | $< 13.1$            | $4.1^{+8.3}_{-3.8}$  | $< 17.3$            |
| $\tau_H \times 10^{-22}$ [s] <sup>†</sup>     | $\mathcal{R}_{\text{gg},VV}^{\text{off-shell}} = 1$ | $7.0^{+82.1}_{-5.5}$ | $> 0.5$             | $1.6^{+23.9}_{-1.1}$ | $> 0.4$             |

\*  $\mathcal{R}_{\text{gg}}^{\text{off-shell}}$  is only weakly constrained and has no defined limits at 1 or 2  $\sigma$  within the scanned range, as seen in [Fig. 11.7](#).

† The Higgs width  $\Gamma_H$  and lifetime  $\tau_H$  are both derived from the  $\mu_{\Gamma_H}$  fit configuration and are calculated using the relation  $\Gamma_H = \Gamma_H^{SM} \cdot \mu_{\Gamma_H}$  and  $\tau_H = \hbar/\Gamma_H$ , respectively.

<sup>1</sup> This 95% CL upper limit is derived using the  $t_\mu$  test statistic, which considers fluctuations in both directions. This will overcover when interpreted as a one-sided limit. The quoted limited limit is therefore somewhat conservative.

Secondly, this thesis introduces two innovative interpretations of the Higgs and electroweak sector into the SM-EFT framework. The first interpretation is an SM-EFT interpretation of STXS Higgs production rates and  $H \rightarrow \gamma\gamma, ZZ^*, WW^*, \tau\tau, b\bar{b}$  decay measurements, in which 13 sensitive directions composed of 31 Wilson coefficients are measured from the data. Although no deviations from the SM are found, the simultaneous fit does provide energy scale limits on the existence of new physics beyond the SM in the Higgs sector.

The second interpretation is an SM-EFT interpretation of the  $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$  and  $pp \rightarrow WW$  processes and illustrates the potential of the SM-EFT interpretation of combined Higgs signal strength and differential electroweak measurements. In this analysis, signal strength and differential leading-lepton  $p_T$  measurements are combined and interpreted in terms of 13 Wilson coefficients and 8 sensitive directions. The results of this interpretation show no deviations from the SM, but do provide a new methodology for the interpretation of combined Higgs and electroweak measurements in the SM-EFT framework.

## Outlook

When looking out to future improvements outside of the scope of this thesis for off-shell, lifetime and SM-EFT measurements, there are several areas that could be improved upon. Many of these new advents, techniques and ideas are discussed in the following sections.

It should be noted that a decade ago, the off-shell Higgs production and width measurements within 100% accuracy were considered to be impossible due to the resolution of the detector and the low cross-section of off-shell Higgs production. Thanks to innovative analysis techniques and the continued collection of data at the LHC, high precision off-shell Higgs production and width measurements are now a reality for ATLAS and CMS. This is a great achievement for the LHC experiments and illustrates the power of the LHC as a tool for precision measurements in the Higgs sector.

## Additional data

As the total off-shell measurement uncertainty is currently statistically limited, the most efficient improvement of the analysis will come from a larger dataset. With the continued collection of high-energy  $pp$ -collision data at the LHC, the prospects for advancing statistically limited analyses like the off-shell analysis and SM-EFT interpretations become increasingly promising. The  $3000 \text{ fb}^{-1}$  dataset expected to be produced by the HL-LHC, will allow measurements of off-shell couplings and SM-EFT parameter space with an unprecedented precision. As the statistical power of the data set grows, it will not only allow for more precise measurements of the off-shell Higgs production cross-section, but likely also improve the data-driven methods used for estimating backgrounds.

Acquiring a larger dataset can also be achieved via the addition of new channels of the off-shell Higgs production, which is the reason the  $ee/\mu\mu$  channel is added to the present

off-shell  $WW$  analysis. The further addition of the semi-leptonic decay channels,  $H^* \rightarrow WW \rightarrow \ell\nu qq$ , would allow for a more complete picture of the off-shell Higgs production and a drastic increase in the number of events available for analysis. Unfortunately, the semi-leptonic decay channels are more challenging to separate from non-interfering background due to the presence of two quarks, a lepton and missing energy in the final state, which is a background signature. However, with the advent of machine learning techniques, it is possible that the separation channels from the background can be improved to overcome this challenge.

### Improved modelling

The ggF signal modelling in the off-shell analysis is based on the SHERPA 2.2.2 generator for matrix element, parton showering and hadronization and interfaced with dedicated higher-order NLO corrections calculated using Powheg-Box. Although experimental, the Powheg-Box matrix element generator can also be matched with Pythia8 for parton showering and hadronization. This would allow for a more consistent treatment of the signal modelling, as the Pythia8 generator is also used for VBF.

Additionally, the direct use of NLO matrix elements would greatly reduce the resummation and matching uncertainties present in the current signal modelling. According to preliminary studies performed by the ATLAS collaboration, the use of a matched Powheg-Box+Pythia8 generator for the ggF signal modelling would reduce the parton shower and hadronization uncertainties by a factor of 2-3. The resummation uncertainty, which is by far the largest uncertainty on the signal, is now of order 50% and would be reduced to 10-20% with the use of a matched generator.

The main bottleneck for the implementation of a matched Powheg-Box+Pythia8 generator is the computational cost of the generator. The parton shower recoil uncertainty would be replaced by a Pythia8 dipole shower uncertainty, which is also expected to be smaller. The calculation of NLO and NNLO  $gg \rightarrow H^* \rightarrow WW$  matrix elements is computationally expensive and requires a large amount of resources. If this bottleneck can be overcome, the use of this new generator would be one of the most impactful improvements for the off-shell analysis. This would not only improve the off-shell  $WW$  modelling, but also the modelling of the off-shell  $ZZ$ , as they share a similar production mechanism.

From all theoretical background uncertainties, the  $t\bar{t}$  background uncertainties are the most limiting for the off-shell  $WW$  measurement. Currently, a 3D reweighing tool is used to improve the  $t\bar{t}$  modelling in the off-shell analysis. However, the reweighing tool is not able to fully correct the mis-modelling of the  $t\bar{t}$  background and therefore comes with a large reweighing uncertainty, which limits further improvement. Furthermore, the use of the SHERPA 2.2.12/2.2.14 generator for the  $Z/\gamma^* + \text{jets}$  background would also benefit from an update to a more recent version of the generator. The SHERPA 2.2.12/2.2.14 generator is known to have issues with modelling the  $p_T$  spectra of the  $Z/\gamma^* + \text{jets}$  background, which is the dominant background in the off-shell  $WW$  same-

flavour analysis. The use of a more recent version of the SHERPA generator would likely improve the modelling of the  $Z/\gamma^* + \text{jets}$  background and reduce the uncertainties on the background estimation in the  $ee/\mu\mu$  channel significantly.

### Improved jet uncertainties

The uncertainties on the energy scale and resolution of the jets are the largest and most impactful reconstruction uncertainties in the off-shell  $WW$  analysis. This is due to the fact that the measurement makes use of the jet multiplicity and the jet  $p_T$  spectra to categorize the signal events into ggF and VBF. Consequently, when the JES and JER uncertainties are large, the categorisation of the events becomes less precise, the signal extraction becomes more difficult and results in the migration of events from one region to another. This, in turn, causes large systematic uncertainty in the yields in each region.

In particular, the uncertainty on the JES pile-up factor  $\rho$  and JES flavour composition are impactful uncertainties in the off-shell  $WW$  analysis. The JES pile-up factor  $\rho$  is conservative and taken as 50%, which could be reduced through the use of a more sophisticated pile-up correction method. The JES flavour composition assumes 50% of all jets are gluon jets, which is an over simplification of the reality. Up and down variations of the JES flavour composition are taken as  $\pm 100\%$  of the nominal value, which is again a conservative assumption. A correction based on a flavour composition estimation of the jet would drop the average uncertainty and improve the categorisation uncertainty of the off-shell signal.

### Off-shell Higgs combinations

One of the most exciting prospects in the off-shell Higgs community is, in my opinion, the combination of the off-shell  $WW$  and  $ZZ$  and  $t\bar{t}H/4\text{-top}$  measurements. As shown in the discussion, the  $WW$  results are competitive with the  $ZZ$  measurement and will likely result in a combined measurement with a significantly higher combined precision. The combination of the off-shell  $WW$  and  $ZZ$  measurements will allow for a more complete picture of the off-shell Higgs production and a more precise measurement of the Higgs couplings. The combination will also allow to compare the Higgs width as measured via the  $WW$  and  $ZZ$  channel, which gives additional information on the model assumptions and the role the different vector bosons play in the off-shell Higgs couplings and width measurements.

While the combination of the off-shell  $WW$  and  $ZZ$  measurements is the most immediate prospect, the combination of the off-shell  $WW$  and  $ZZ$  measurements with the off-shell  $t\bar{t}H$  and 4-top measurements have also been a recent point of interest. The off-shell  $t\bar{t}H$  and 4-top constraints on the Higgs total width have recently been published [265] and use a combination of on-shell  $t\bar{t}H$  Higgs boson production and off-shell 4-top production to constrain the total Higgs width. This combination arrives at an expected upper limit of 450 MeV at 95% CL. While not the most precise measurement of the Higgs width, the



measurement is the first to use the Higgs-top Yukawa coupling to measure the Higgs width and is a promising first step towards a more complete picture of the Higgs width and the model assumptions taken to measure it.

A combination of the ATLAS off-shell  $WW$  and  $ZZ$  measurements with the off-shell  $t\bar{t}H$  and 4-top measurements would allow for a very sensitive measurement of the off-shell couplings and Higgs width and may potentially provide a first hint of new physics beyond the SM in the Higgs sector, since the combined measurement approaches the statistical power in order to claim  $5\sigma$  discoveries. When this measurement is combined with the CMS off-shell  $WW$  and  $ZZ$  measurements, the statistical power of the combined measurement becomes even higher and the sensitivity towards new physics beyond the SM therefore is correspondingly larger.

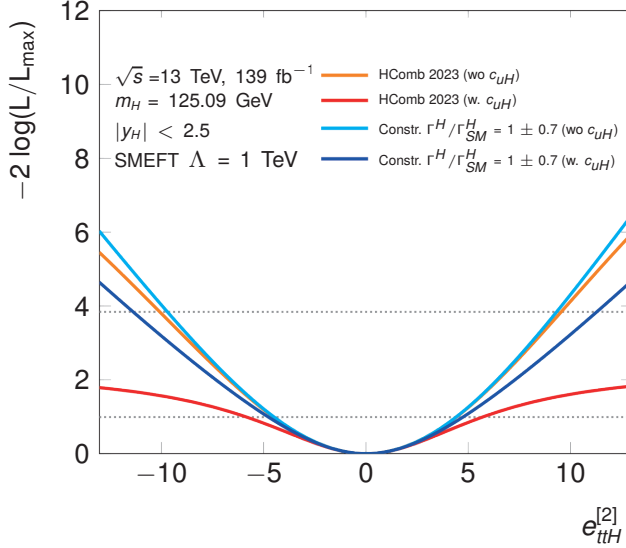
### The next steps for SM-EFT interpretations

A clear way forward for SM-EFT interpretations is to keep adding more measurements from different sectors to a global combination. This is possible since the SM-EFT effects are not limited to the Higgs or EW sector. A good example of this is the addition of top measurements in the SM-EFT interpretations towards the Higgs sector, as  $t\bar{t}$  and  $Wt$  are a major background to Higgs measurements.

A second way of adding more information to the SM-EFT global fits is the addition of CP-odd sensitive variables and processes. The SM-EFT interpretations presented in this thesis do not have sensitivity towards CP-odd variables since the underlying discriminant variables used are CP-even. Current developments are already undergoing for a  $H^* \rightarrow WW \rightarrow \ell\nu\ell\nu$   $\Delta\phi_{jj}$  analysis, which reconstructs the angle between the two jets originating from the ggF +2jet and VBF Higgs production mechanisms. It makes use of an STXS-like binning of  $\Delta\phi_{jj}$  and has good sensitivity towards the  $\mathcal{O}_{H\tilde{W}}$ ,  $\mathcal{O}_{H\tilde{W}B}$  and  $\mathcal{O}_{H\tilde{B}}$  operators through VBF production and also the  $\mathcal{O}_{H\tilde{G}}$  operator through ggF production.

Another possible direction for the SM-EFT interpretations in the Higgs sector is to include the off-shell Higgs measurements in the combined fits, as proposed by the off-shell interpretations task force of the LHC Higgs Working Group [266]. A first step in this direction has already been taken by the ATLAS collaboration, where the off-shell  $ZZ$  measurement has been interpreted in terms of the  $\mathcal{O}_{HG}$  and  $\mathcal{O}_{tH}$  operators [267]. The benefit of including the off-shell measurements in the global fits is that the off-shell measurements are sensitive to the Higgs couplings in a different kinematic regime. This allows for a more complete picture of the Higgs couplings and a better constraint on the SM-EFT parameter space where the on-shell measurements are not sensitive. It also reduces correlations between parameters in the global fits, as the operators shared between the on-shell and off-shell measurements have different effects on the two measurements.

It was shown in the 2022 Higgs workshop in Tokyo that the SM-EFT parametrization of the Higgs sector can benefit from a constraint on the total Higgs width, which can be achieved via the inclusion of an off-shell measurement. An example of this is shown in Fig. 15.1, where the expected likelihood scan of the eigenvector  $e_{ttH}^{[2]}$  is shown as a function of four fit/width constraint configurations, using the statistical model of the Higgs interpretation paper [268].



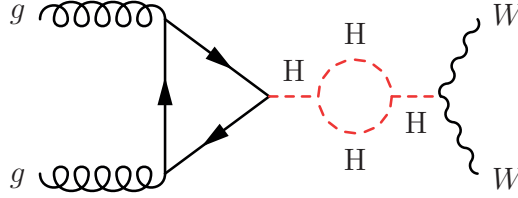
**Figure 15.1:** The impact on the likelihood of the eigenvector  $e_{ttH}^{[2]}$  when including and excluding a constraint on the total Higgs width in the global fit. The likelihood is shown as a function of the  $e_{ttH}^{[2]}$  parameter. The red line shows the likelihood of the  $e_{ttH}^{[2]}$  parameter with an unconstrained width and the inclusion of  $\mathcal{O}_{uH}$ , the blue line shows the improved likelihood when a total Higgs width measurement introduces a  $\Gamma_H/\Gamma_H^{\text{SM}} = 1 \pm 0.7$  Gaussian constraint. The cyan and orange lines show the likelihood when the  $\mathcal{O}_{uH}$  operator is not included, with and without the width constraint.

When the total Higgs width is left unconstrained via the inclusion of the  $\mathcal{O}_{uH}$  operator, due to a large measurement uncertainty on the corresponding Wilson coefficient, the likelihood of the  $e_{ttH}^{[2]}$  parameter, shown in red, becomes flatter, and does not rise to the 95% CL threshold level for any value of the parameter. However, when the total Higgs width is constrained, as shown in the model represented by the blue line, which introduces a comparable sensitivity as obtained from the off-shell  $H \rightarrow WW$  width measurement, the original sensitivity of the  $e_{ttH}^{[2]}$  parameter is restored. Even when the  $\mathcal{O}_{uH}$  operator and its dependence on the total Higgs width are not included, a slight increase in sensitivity is observed, as shown in the model represented by the light blue line.

Furthermore, the off-shell rate measurements can provide sensitivity towards trilinear Higgs couplings [269]. By having access to this coupling, an SM-EFT parameterization of the off-shell measurements can provide limits on  $\lambda$ , which in part determines shape of the Higgs potential:

$$V_{SM}^H = \frac{1}{2}m_H^2 H^2 - \lambda v H^3 + \frac{1}{4}\lambda H^4.$$

An example of a trilinear Higgs coupling is shown in Fig. 15.2, where a Feynman diagram of a Higgs coupling to itself is shown. The diagram can lead to a  $\mathcal{O}(\lambda)$  correction to



**Figure 15.2:** A Feynman diagram of Higgs self-coupling in the SM, which can lead to a  $\mathcal{O}(\lambda)$  correction to the  $gg \rightarrow H^* \rightarrow WW$  process.

the  $gg \rightarrow H^* \rightarrow WW$  process, which can potentially be measured in the off-shell  $WW$  analysis because of the high invariant mass the analysis targets. Via this, and two other diagrams, the process is sensitive to the  $\mathcal{O}_6$  and  $\mathcal{O}_H$  operator. The Wilson coefficients of these operators can be related to a parameter which linearly scales the trilinear coupling ( $-\lambda c_3 v H^3$ ), i.e.  $c_3 = 1 + c_6 - 3/2c_H$ . By using this interpretation method, complementary results to the Higgs self-coupling strength and Higgs potential can be obtained.

At last, a more subtle advent of the implementation of SM-EFT, into width measurements in particular, is the model-independent nature of the effective field theory. Since the SM-EFT parametrization allows for a model-independent interpretation of the Higgs couplings across the full mass range of the Higgs boson, the assumptions currently made when measuring the Higgs width can be relaxed.

At the moment, when performing the Higgs width measurement, the assumption is made that the Higgs boson is a SM Higgs boson and the Higgs couplings are SM-like and cancel out. When the Higgs width is measured in the context of a SM-EFT parametrization, it can be measured without the assumption of the SM Higgs couplings. This will result in a more model-independent measurement of the Higgs width.

## NEDERLANDSE SAMENVATTING

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Het geven van antwoorden op diepgaande vragen over materie en de wetten van de natuur; *dát* is het doel van deeltjesfysica. Dit is deels bereikt door de subatomaire deeltjes om ons heen en de krachten die daartussen spelen te beschrijven. Het standaard model (SM) is het resultaat van deze inspanning en biedt een theoretisch model waarin alle ontdekte deeltjes zijn opgenomen. Echter, met de deeltjes die tot nu toe zijn aangetroffen - beginnend met het elektron ( $e$ ) en eindigend bij het Higgsboson ( $H$ ) - kan het SM slechts een deel van het universum en drie van de vier fundamentele krachten verklaren. Ook al is er nog geen direct bewijs voor nieuwe fysica, het SM wordt door deze tekortkomingen over het algemeen als onvolledig beschouwd.

De ontdekking van het Higgsboson met de Large Hadron Collider<sup>1</sup> (LHC) stelt experimenteel natuurkundigen sinds 2012 in staat om, op een nieuwe manier, naar deeltjes buiten het SM te zoeken. Dit nieuwe hulpmiddel is: het Higgsboson zelf. Dit proefschrift draagt bij aan deze zoektocht door het Higgsboson te bestuderen door middel van het verval in vier leptonen ( $\ell$ ): bestaande uit elektronen, muonen ( $\mu$ ) en neutrino's ( $\nu$ ), via twee geladen vectorbosonen ( $W^\pm$ ). Dit proces wordt veroorzaakt door de zwakke kernkracht, dat bijvoorbeeld ook leidt tot het  $\beta^-$  verval in onstabiele neutron-rijke atomen.

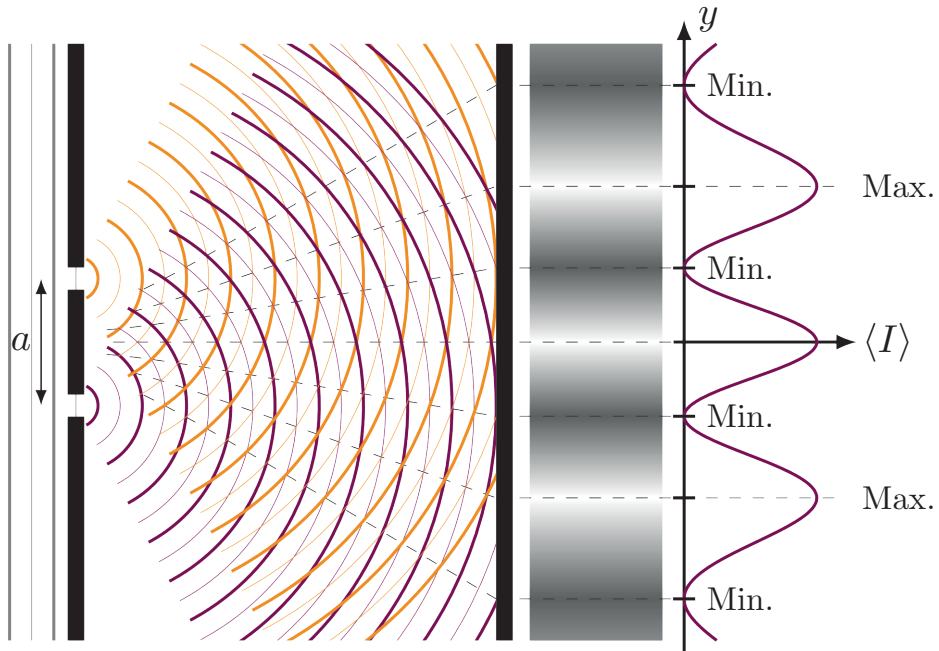
In deze thesis zijn twee verschillende analysetechnieken gebruikt om een beter begrip te krijgen van het Higgsboson. Beide maken gebruik van de interferentiepatronen tussen het Higgsboson en andere deeltjes in het SM. De interferentiepatronen worden vervolgens ingezet om eigenschappen van het Higgsboson te meten, of om naar nieuwe interacties buiten het SM te zoeken. Deze interferentie is een kwantummechanisch effect dat optreedt omdat subatomaire deeltjes een golfkarakter vertonen.

Thomas Young observeerde dit effect voor het eerst in 1801 door een coherente lichtstraal te schijnen op twee dunne lichtdoorlatende openingen. In het beroemde tweespleten-experiment (Figuur 16.1) interfereerden de fotonen ( $\gamma$ ) in de twee lichtstralen door een faseverschil.

De afstand  $a$  tussen de twee spleten zorgt voor een padlengteverschil, waardoor de fotonen zowel constructieve als destructieve interferentie ondergaan. Dit leidt tot afwisselend lichte en donkere banden op het scherm, waarmee niet alleen het golfkarakter van het licht werd aangetoond, maar ook die van de  $\gamma$ -deeltjes.

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<sup>1</sup> De LHC bevindt zich onder het CERN-instituut en het kanton Genève, op de grens tussen Frankrijk en Zwitserland.



**Figuur 16.1:** Een schematische weergave van het experiment van Thomas Young dat resulteert in een interferentiepatroon tussen de oranje en paarse lichtgolf geprojecteerd op een donker scherm door een faseverschil. In plaats van een uniforme lichtverdeling, veroorzaakt het interferentie-effect afwisselende lichte en donkere banden omdat de fotonen zich als een golf gedragen [12].

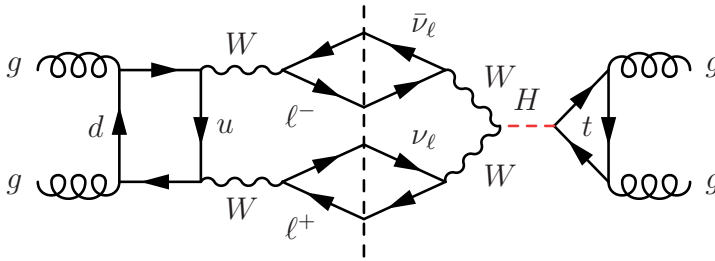
## Off-shell Higgsboson en levensduur metingen

Metingen van het Higgsboson in het off-shell regime (buiten de resonantie-massa) bieden de mogelijkheid om de minder toegankelijke aspecten van dit deeltje te onderzoeken. Het stelt ons namelijk in staat om zijn gemiddelde leeftijd en zijn koppeling sterkte te meten, zonder al te veel aannames. Dit kan worden beïnvloed door onbekende deeltjes die koppelen met het Higgsboson op een massa-schaal hoger dan zijn on-shell resonantie-massa  $m_H = 125 \text{ GeV}$ .

De data verzameld door het ATLAS experiment tijdens Run-2 van de LHC in de jaren 2015-2018, bij een zwaartepuntsenergie van  $13 \text{ TeV}$ , zijn gebruikt voor deze meting. De totale dataset komt overeen met  $140 \text{ fb}^{-1}$ , wat een maat is voor het aantal proton-proton botsingen dat plaatsvond tijdens de dataverzameling en is de grootste hoeveelheid data ooit gebruikt voor dit type analyse.

De off-shell analyse meet het  $pp \rightarrow H \rightarrow WW \rightarrow \ell^+ \nu \ell^- \bar{\nu}$  signaal in het off-shell continuüm. Hierin heeft het Higgsboson zelf een hoge massa, die zo hoog is zijn productiesnelheid niet langer afhangt van zijn levensduur. Dit gebeurt ongeveer 15% van alle keren dat een Higgsboson is geproduceerd, en is relatief groot vanwege een tweede en derde resonantie bij de energieën  $m_{WW} = 161 \text{ GeV}$  en  $m_{t\bar{t}} = 346 \text{ GeV}$ . In het off-shell

continuüm interfereert dit signaal eveneens sterk met het proces  $gg \rightarrow WW \rightarrow \ell^+ \nu \ell^- \bar{\nu}$ , dat een identieke begin- en eindtoestand heeft. De productiemodus die bij deze drie resonanties hoort, wordt gluon-gluon fusie (ggF) genoemd. Het Feynman diagram van deze productie, samen met zijn interfererende achtergrond, worden getoond in [Figuur 16.2](#). Tussen deze twee processen treedt sterke interferentie op. Deze is negatief vanwege het faseverschil, welke veroorzaakt wordt door de verschillende koppelingstructuur die wordt gebruikt om dezelfde eindtoestand te bereiken.

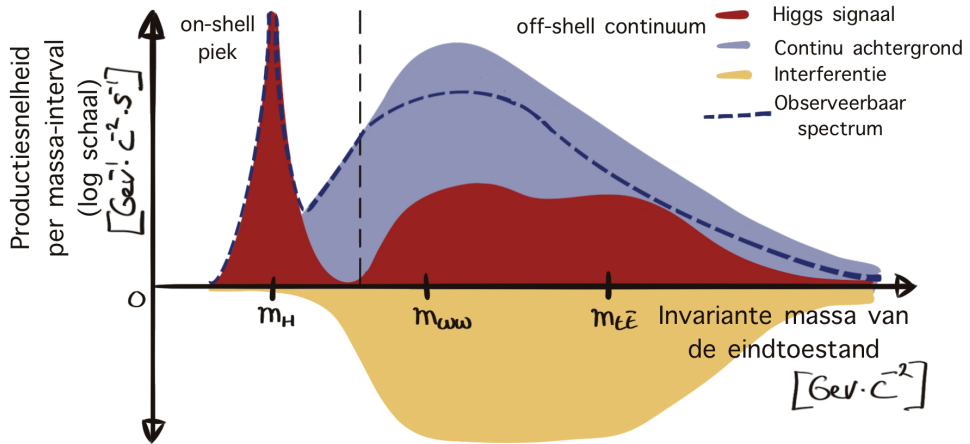


**Figuur 16.2:** Twee interfererende Feynman diagrammen; het Higgsboson (rode stippellijn) wordt rechts geproduceerd via botsende gluonen ( $g$ ) en een driehoekus met topquarks ( $t$ ), en valt uiteen in  $\ell^+ \nu \ell^- \bar{\nu}$  via twee  $W^\pm$ -bosonen. Links staat het achtergrondproces dat hiermee interfereert. Dit is duidelijk gemaakt door het diagram aan de rechterkant gespiegeld weer te geven langs de zwarte stippellijn.

De ggF-productiemodus produceert een Higgsboson door de  $t\bar{t}$ -resonantie, geïnitieerd door twee gluonen, en vervalst via een  $WW$ -resonantie. [Figuur 16.3](#) toont een schets van de resulterende on-shell en off-shell Higgs massa-distributie, samen met de continuüm achtergrond en het negatieve interferentie-effect tussen de twee processen. Dit is vergelijkbaar met het interferentie-experiment van Thomas Young en resulteert in een interferentiepatroon tussen het Higgsboson en de continuüm achtergrond in het waarneembare spectrum. De som van alle bijdragen vormt dit waarneembare spectrum en is vervolgens vergeleken met simulaties om tot een meting van de off-shell Higgsboson productiesnelheid te komen.

Een vergelijkbaar spectrum wordt geproduceerd via vectorboson-fusie (VBF) en is tevens in deze analyse gemeten. Dit signaal produceert Higgsbosonen via quark-verstrooiing en vectorboson ( $V = W^\pm, Z$ ) interacties. Dit proces produceert ook twee quarks en, hoewel niet onderhevig aan extra resonanties die de productie verhogen, bleek door hoge gevoeligheid een waardevolle toevoeging aan de analyse.

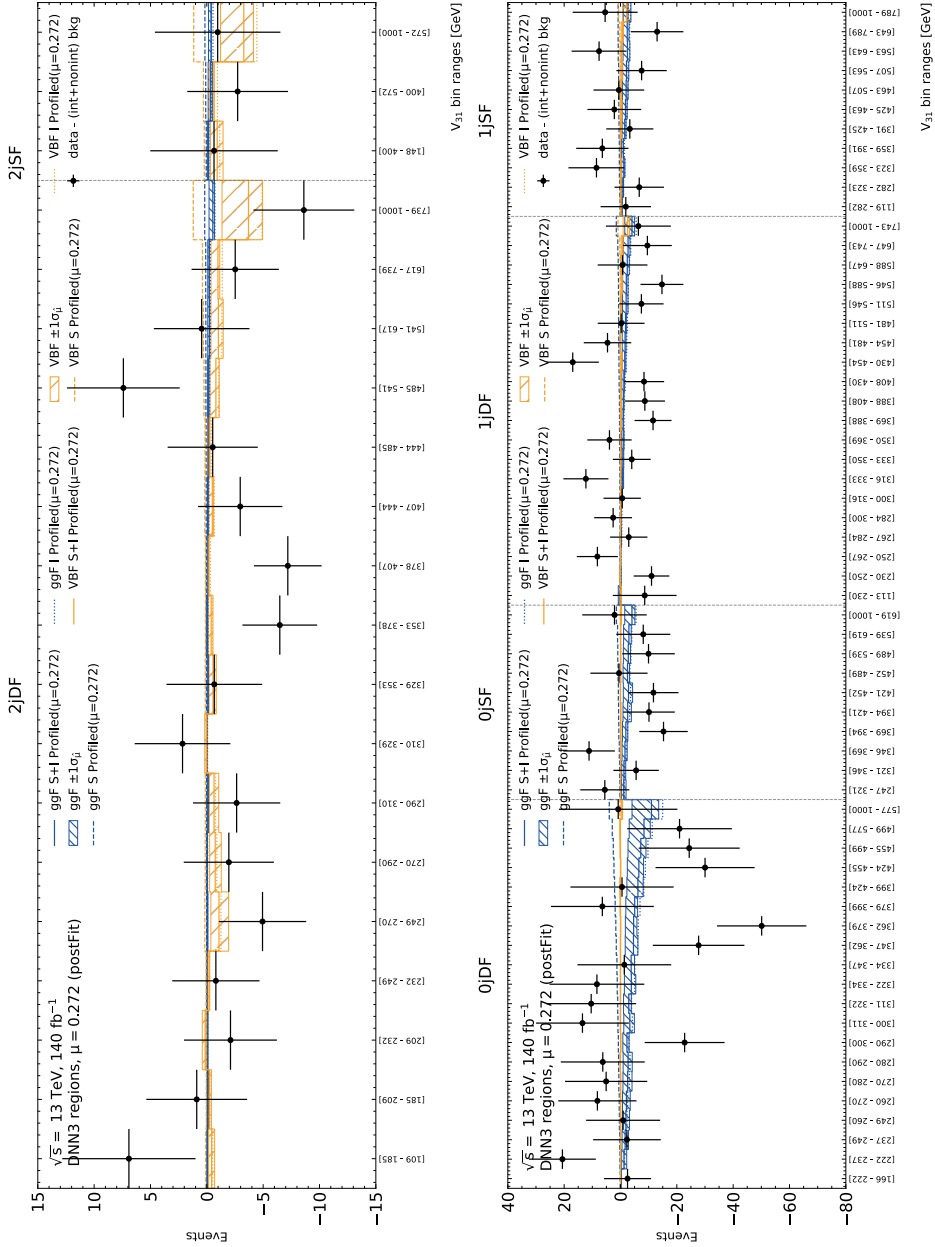
Om de signaalsterkte van dit proces te meten, past de off-shell analyse eerst een voorselectie toe om oninteressante proton-proton botsingen te verwijderen. Hierin wordt de aanwezigheid van twee leptonen boven een  $p_T$ -drempel en een gecombineerde minimale invariante massa vereist. Daarnaast worden botsingen verdeeld aan de hand van hun verval producten over de  $ee/\mu\mu$  en  $e\mu/\mu e$  categorieën. Na de voorselectie worden botsingen ingedeeld op basis van jet-activiteit in de  $N_{\text{jet}} = 0$ ,  $N_{\text{jet}} = 1$  en  $N_{\text{jet}} \geq 2$ -categorieën.



**Figuur 16.3:** Een schets van de on-shell resonantie rond de Higgs massa ( $m_H$ ) en het off-shell continuüm waarin de  $WW$  en  $t\bar{t}$  resonantie aanwezig zijn, op de  $m_{WW}$  en  $m_{t\bar{t}}$  massa's. Het Higgs-signaal en de interfererende continuüm achtergrond zijn in rood en paars getoond. Het interferentie-effect is in geel weergegeven en het totale waarneembare spectrum is met de donkerpaarse gestreepte lijn weergegeven.

Hierna zijn een reeks orthogonaliteit- en topologie-selecties gemaakt om het signaal verder te versterken ten opzichte van de achtergrond. De laatste stap is het verdelen van zes opgebouwde signaalregio's in drie subsignaalregio's door het gebruik van de  $V_{0.3|1.0}$  variabele (sterk gecorreleerd met  $m_{WW}$ ) en zes diepe neurale netwerken (DNN). De DNNs zijn getraind om interfererende achtergrond en signaal te scheiden van niet-interfererende achtergronden en zo de gevoeligheid van de analyse te verbeteren.

In [Figuur 16.4](#) zijn de gereconstrueerde ggF- en VBF-signalen in de meest gevoelige regio's van de analyse getoond. Hierin is de gemeten achtergrond afgetrokken van de data, waardoor alleen signalen overblijven. Daarnaast zijn de gesimuleerde signaal- en interferentie-componenten getoond. Deze visualisatie stelt de lezer in staat om een gevoel te krijgen voor het gemeten interferentie-effect en signaal van het Higgsboson, in het onderzochte off-shell regime. De data komen overeen met het gemodelleerde signaal-plus-interferentie effect (S+I), aangegeven door de doorgetrokken lijnen, en laten het volledige effect van het off-shell Higgsboson zien. De individuele (positieve) signaal en grotere (negatieve) interferentie componenten zijn getoond door middel van de stippel- en gestreepte lijnen. Het gecombineerde signaal-plus-interferentie effect, dat het off-shell Higgsboson heeft op de distributie, is in deze regio gemeten als een algemeen destructief interferentie-effect. De meting gaf een lagere signaalsterkte dan verwacht, deels door een gebrek aan waargenomen negatieve interferentie. Het resultaat is daarentegen nog steeds in overeenstemming met de voorspelling van het SM.



**Figuur 16.4:** De gereconstrueerde signaal data van de ggF (blauw) en VBF (oranje) off-shell meting, waarbij alleen de signaal en interfererende (S+I) bijdragen overblijven. De doorgetrokken lijn vertegenwoordigt het S+I effect veroorzaakt door het off-shell Higgsboson. De stippellijn toont alleen het off-shell signaal en de gestreepte lijn illustreert alleen de interferentie. De onzekerheid in het bovenste en ratio-gedeelte komt overeen met de totale post-fit theoretische, experimentele en statistische onzekerheid.



Het gemeten signaal is vervolgens vergeleken met de achtergrondhypothese om de significantie te berekenen. Hierbij kwam de off-shell signaalsterkte van  $\mu_{\text{off-shell}} = 0.3_{-0.3}^{+0.9}$  overeen met een waargenomen (verwachte) significantie  $Z$  van  $0.7\sigma$  ( $1.2\sigma$ ). Dit resultaat stelt de off-shell analyse in staat om een limiet van 3.4 keer het SM, met een betrouwbaarheidsniveau van 95%, vast te stellen voor de off-shell koppelingssterkte.

Bovenop dit resultaat is de off-shell analyse ook gecombineerd met de on-shell  $H \rightarrow WW^*$  analyse, met de vergelijking:

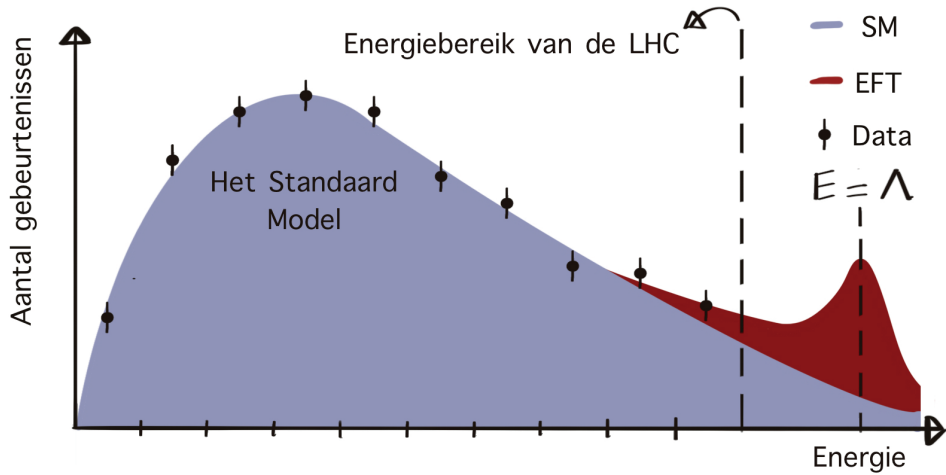
$$\mu_{\text{off-shell}} = \mu_{\text{on-shell}} \cdot \frac{\Gamma_H}{\Gamma_H^{\text{SM}}}.$$

Deze combinatie is uitgevoerd om de Higgsboson resonantie-breedte en gemiddelde leeftijd te meten. Het resultaat van de resonantie-breedte meting is  $\Gamma_H = 0.9_{-0.9}^{+3.4}$  MeV. Dit komt overeen met een gemiddelde leeftijd van  $\tau_H = 7.0_{-5.5}^{+82.1} \times 10^{-22}$  s en een bovengrens voor de totale breedte van het Higgsboson van  $\Gamma_H < 13.1$  MeV, met een betrouwbaarheidsniveau van 95%. Dit resultaat komt overeen met de SM Higgsboson voorspelde breedte van 4.1 MeV. Er is tevens een bovengrens vastgesteld voor de off-shell/on-shell koppeling-ratio  $R_{VV}$  van 2.8, met een betrouwbaarheidsniveau van 95%. Deze analyse resulteerde in de beste limiet op de breedte van het Higgsboson met behulp van het  $H \rightarrow WW$  vervalskanaal tot nu toe, en heeft zich bewezen ten opzichte van vergelijkbare metingen met  $H \rightarrow ZZ$  vervalstopologieën.

### Interpretaties van het standaard model via effectieve veldentheorieën

De interpretaties van Higgs metingen in dit proefschrift hebben als doel operatoren en koppelingen te ontdekken die kunnen duiden op fysica niet beschreven in het SM. Deze kunnen in de effectieve veldentheorie worden veroorzaakt door een nieuw zwaar deeltje op een energieniveau van  $\Lambda = 1$  TeV. In dit proefschrift is de meting van deze koppelingen bereikt door de interpretatie van twee gecombineerde metingen van de Higgs- en zwakke kernkracht sectoren als functie van de SM effectieve veldentheorie (in het Engels afgekort als SM-EFT). Deze interpretatie is vervolgens gebruikt om te zoeken naar afwijkingen van de SM-voorspellingen. Dit kan bijvoorbeeld worden veroorzaakt door nieuwe fysica die niet is opgenomen in het SM, maar die wel aanwezig is op hogere energieniveaus buiten het bereik van de LHC, zoals geschetst in [Figuur 16.5](#).

De eerste interpretatie past de SM-EFT toe op Simplified Template Cross-Section (STXS) metingen van Higgs-productie en het verval in  $\gamma\gamma$ ,  $ZZ^*$ ,  $WW^*$ ,  $\tau\tau$  en  $b\bar{b}$  deeltjes. De interpretatie gebruikt data van proton-proton botsingen, verzameld door het ATLAS experiment tijdens Run-2 van de LHC in 2015-2018 bij  $\sqrt{s} = 13$  TeV, en komt overeen met  $139 \text{ fb}^{-1}$ . Deze analyse biedt een groot aantal signaalsterkte metingen van verschillende delen van de transversale impuls distributie van de Higgs ( $p_T^H$ ), de invariante massa van twee jets ( $m_{jj}$ ) en de transversale impuls van het bijbehorende vectorboson ( $p_T^V$ ). De gemeten signaalsterktes, als functie van kinematica van het Higgsboson productie-proces, geven een hoge gevoeligheid voor een groot aantal operatoren die verband houden met



**Figuur 16.5:** Een schets van mogelijke afwijkingen van het SM als functie van het energieniveau. Het SM (paars) kan de data alleen beschrijven bij lage energieniveaus. Het SM-EFT model, met een extra toevoeging van nieuwe fysica (rood), is nodig om de data bij hogere energieniveaus te beschrijven.

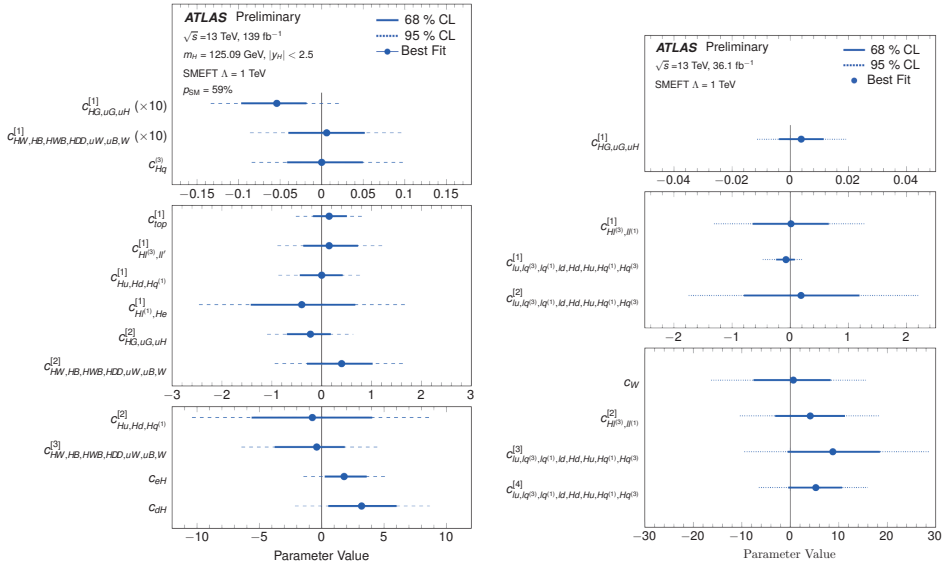
de Higgs sector. Dit geeft een zeer gedetailleerde beeld van het Higgsboson in termen van SM-EFT operatoren, wat verder is versterkt door de combinatie van verschillende productiekanalen en vervalprocessen.

Het grote aantal operatoren, die de productie en het verval van het Higgsboson beïnvloeden, kunnen niet tegelijk gemeten worden vanwege de beperkte experimentele gevoeligheid. Daarom is een subset van gevoelige parameters gedefinieerd door een eigenvector-analyse met de gemeten data. De nieuwe gedefinieerde richtingen zijn lineaire combinaties van SM-EFT operatoren en worden opgebouwd door rekening te houden met de beschikbare fysische informatie en de verwachte gevoeligheid van de analyse. Met deze techniek worden 13 richtingen, bestaand uit 31 Wilson coëfficiënten, tegelijk gemeten. Hierdoor zijn in de Higgs sector de gemeten waarden en limieten van SM-EFT parameters in kaart gebracht.

De tweede interpretatie in dit proefschrift analyseert de  $pp \rightarrow H \rightarrow WW$  en de  $pp \rightarrow WW$  processen door het oog van de SM-EFT en benadrukt de waarde van het combineren van metingen van verschillende natuurkunde-sectoren om SM-EFT operatoren te meten. De gecombineerde analyses betreffen een signaalsterkte-meting en een differentiële lepton- $p_T$  meting van het meest energetische lepton. Beide maken gebruik van een gedeeltelijke Run-2 dataset, verzameld door het ATLAS experiment in 2015-2016 met  $\sqrt{s} = 13$  TeV, en komen overeen met  $36.1 \text{ fb}^{-1}$ .

In de interpretatie zijn de twee verschillende meetmethoden gecombineerd met behulp van een nieuw coherent statistisch kader dat acht gevoelige richtingen, geconstrueerd via een eigenvector-analyse, uit 13 Wilson coëfficiënten. De resultaten van de twee

interpretaties zijn naast elkaar weergegeven in [Figuur 16.6](#): deze tonen geen afwijkingen binnen  $2\sigma$  wanneer ze worden vergeleken met het SM. De metingen beperken wel het bestaan van nieuwe fysica op een schaal van 1 TeV, en kunnen worden geëxtrapoleerd naar hogere energie-niveaus wanneer nodig. Met deze nieuwe methoden kunnen SM-EFT interpretaties breed worden toegepast, inclusief off-shell Higgs- en breedtemetingen, waardoor een preciezer en model-onafhankelijker beeld van de Higgs sector ontstaat.



**Figuur 16.6:** De samengevatte Wilson-coëfficiënt resultaten. De getoonde limieten komen overeen met 68% (doorgetrokken) en 95% (gestreept) betrouwbaarheidsintervallen, waarbij alle andere gevoelige richtingen tegelijkertijd worden gemeten. De linker grafiek toont de resultaten van de gecombineerde Higgs metingen, terwijl de rechter grafiek de resultaten van de gecombineerde  $H \rightarrow WW^*$  en  $WW$  metingen toont.

Dit proefschrift heeft grenzen gesteld aan nieuwe fysica en koppeling die zich verbergen binnen het hoge energieregime. Ook heeft het met de meting van het off-shell Higgs-boson en de Higgs-levensduur meer inzicht gegeven in de potentiële deeltjes die de Higgsboson-vervaltijd kunnen beïnvloeden. Samen vormen deze resultaten een alomvattende zoektocht naar nieuwe fysica in de Higgs sector en geven ze een completer beeld van het Higgsboson en zijn eigenschappen.

Als ik terugkijk op mijn jaren in het veld, realiseer ik me dat het Higgs-precisietijdperk nu officieel is begonnen. De komende jaren komen er veel interessante resultaten en ontdekkingen - houd dit dus goed in de gaten!

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## ABOUT THE AUTHOR

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Bryan Kortman received his Bachelor of Science in Applied Physics in 2016 at the University of Twente. He specialized in material physics and conducted his studies within the Energy, Materials, and Systems (EMS) group, at the faculty of science and technology. During his time in the EMS group, he completed an internship on the design and fabrication of high-temperature superconducting solenoids at CERN, within the ATLAS magnet research and development group. After this, he took extracurricular courses from the gravitation, astro-, and particle physics (GRAPPA) track at the University of Amsterdam. In Amsterdam, he completed his master's thesis in 2019 at the Dutch National Institute for Subatomic Physics (Nikhef).



Since December 2019, he has been pursuing his PhD in particle physics at the University of Twente, with his research activities based at Nikhef in Amsterdam and CERN in Geneva. He was involved in the ATLAS Level-1 missing transverse energy group, where he worked on using offline event reconstruction, provided by the reconstruction framework Athena, to improve low latency online triggers through Kalman filters and differential programming. During his PhD, he interpreted the 2022 ATLAS combined Higgs production and decay measurements, as well as the combined  $H \rightarrow WW^*$  and  $pp \rightarrow WW$  results from 2016, within the SM-EFT framework, deriving upper limits for Wilson coefficients and exploring physics beyond the Standard Model. Additionally, he performed the ATLAS Run-2  $H^* \rightarrow WW$  off-shell measurement and combined it with the on-shell  $H \rightarrow WW^*$  analysis to obtain a measurement and limit on the total Higgs width and its lifetime.

While working on his PhD, he also performed research on the performance of typical artificial intelligence workloads in high-energy physics and condensed matter physics on various computing architectures (CPU, GPU, in-memory computing, and FPGAs). This work was carried out on a part-time basis as part of the Green IT team, a collaboration with Radboud University, IBM, and SURF.









