

MEASUREMENT OF ELECTROWEAK WZ BOSON
PRODUCTION AND SEARCH FOR NEW PHYSICS IN
PROTON-PROTON COLLISIONS AT $\sqrt{s} = 13 \text{ TeV}$ WITH
THE CMS DETECTOR AT THE CERN LHC

by

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Abstract

A measurement of the simultaneous production of a W and a Z boson (WZ production) via vector boson scattering is presented. The measurement is performed in the leptonic decay modes $WZ \rightarrow \ell\nu\ell'\ell'$, where ℓ, ℓ' indicate an electron or muon. The analysis is based on a data sample of proton-proton collisions at $\sqrt{s} = 13$ TeV at the CERN Large Hadron Collider collected with the Compact Muon Solenoid detector and corresponding to an integrated luminosity of 35.9 fb^{-1} . The WZ plus two jet production cross section is measured in fiducial regions with enhanced contributions from electroweak (EW) production and found to be consistent with standard model predictions. The EW WZ production in association with two jets is measured with an observed (expected) significance of 2.2 (2.5) standard deviations. Results are also interpreted in terms of physics beyond the standard model modifying the interactions of the W and Z bosons. Constraints on charged Higgs boson production and on anomalous quartic gauge couplings in terms of dimension-eight effective field theory operators are presented.

This thesis is dedicated to the memory of Dr. David Murdock. Your passion for particle physics spread to me, and your generosity with your time and knowledge gave me and many other students the foundation to be successful beyond the walls of Tennessee Technological University.

“Les orages, la brume, la neige, quelquefois ça t’embêtera. Pense alors à tous ceux qui ont connu ça avant toi, et dis-toi simplement : ce que d’autres ont réussi, on peut toujours le réussir.”

— Antoine de Saint-Exupéry, *Terre des Hommes*

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Chapter 1

Introduction

1.1 The universe of particles

The ambitious goal of particle physics is to identify a set of fundamental building blocks from which all observed phenomena of the universe can be constructed. A final description of the fundamental, indivisible “particles” that form the matter of our universe should be both minimal and complete; it should provide a set of particles and interactions capable of describing all observed phenomena, but it should always strive to reduce the set of basic elements if possible. Perhaps nature did not need to realize our human desire for such simplicity, but this paradigm has been remarkably successful [1]. Though holes and unanswered questions remain, interactions of a countable number of particles accurately describe a vast array of phenomena from microscopic to cosmological distance scales.

Identifying the fundamental particles and their interactions is an endeavor that necessitates constant validation and refinement. For over a century, particle physics has been characterized by experimental discoveries, many predicted by theoretical inferences, and others completely unexpected. They have dramatically propelled the

field forward, while at times upending the accepted thought. The history of particle physics is a fascinating one, with a subtle and intricate narrative justifying a much deeper discussion than presented here. Ref. [1] gives a technical introduction to the field in a historical context, whereas Ref. [2] gives a purely historical account. An encyclopedic overview of the major theoretical and experimental advances in particle physics is given in Ref. [3].

The modern genesis of the field can reasonably be traced to the discovery of the electron in 1897 by J. J. Thomson [4], a particle still believed to be elementary today. At the time, though supported by the thermodynamic studies of Boltzman, the once-popular idea of a quantized unit of matter was held in widespread doubt. The investigation of radioactive decays, by Becquerel, Marie and Pierre Curie, Rutherford, and others [5], became closely linked with efforts to categorize the structure of the chemical elements, assumed to be building blocks of nature. However, as the atomic model was accepted, it quickly became clear that the atom itself is not elementary. At the turn of the century, the electron was incorporated into new models of a composite atom, and the work of Rutherford demonstrated the presence of the positively charged nucleus at the center of the atom [6]. The positively-charged proton, with a mass nearly two thousand times that of the electron, was discovered in 1919, again by Rutherford, which he correctly associated with the nucleus [7].

A mathematical understanding of the electromagnetic interactions governing the proton and electron had already been established by James Clerk Maxwell decades before their discoveries. Combined with the new quantum theory, rapidly developing in the 1920s and 1930s, the electromagnetic interaction is also sufficient to explain the stability and arrangements of electrons into atomic orbitals around the nucleus. A key aspect of this model is the fact that the electron, and all other fundamental particles, carry a fundamental angular momentum and magnetic moment, termed

“spin.” As required by the quantum theory, and demonstrated by Stern and Gerlach in 1922 [8], spin is quantized. The electron and proton spins are always observed to be $\pm\hbar/2$, where $\hbar = 1.055 \times 10^{-34} \text{ J} \cdot \text{s}$ is the reduced Planck constant. We refer to such half-integer spin particles as *fermions*. Particles with scalar spin are referred to as *bosons*. In particular, spin-1/2 objects are known as spinors, whereas spin-1 objects are vectors.

A major success of the quantum theory is the resolution of an apparent conflict between the wave-like and particle-like nature of light: in relativistic quantum mechanics (quantum field theory, or QFT), all particles are associated with fields, giving them both wave-like and particle-like properties. By the early 1940s, a relativistic quantum theory of the charged particles and their interactions via the electromagnetic force had been developed in this language. Quantum electrodynamics (QED) is a QFT in which interactions follow from the invariance of a spin-1 field under a class of transformation of the field properties known as *gauge transformations*. In such a quantum gauge theory, the *gauge symmetry* of the spin-1 field leads to interactions between the fermions and boson fields. In the theory of quantum electrodynamics (QED), electromagnetic interactions between fermions are communicated by the exchange of the spin-1 photon (γ), the particle of light. Performing calculations in QED presents significant challenges, but as mathematical techniques were developed to extract meaningful results, the predictive power of QED was astounding.

Yet problems with the atomic model remained. The mass of the atom suggested another, neutral component of the nucleus. The neutron, with nearly equal mass to the proton, was discovered by Chadwick in 1932 [9, 10]. A new force, the “strong force” was proposed as a means to confine the neutral and positively charged constituents of the nucleus, but its mechanism was completely unknown. Further confusion arose from atomic decays. Whereas many atomic decays are consistent with

a breakdown of the original atom into smaller atomic pieces, β decays, in which a radioactive element with atomic number A (the number of protons) decays to an element with $A + 1$ via the emission of an electron, are not.

The first predictive theoretical treatment of β decay was developed by Fermi in 1934 [11]. At a time when the structure of the nucleus was still not convincingly established, he recognized that all β decays could be generalized to the process $n \rightarrow p + e$ (or $n \rightarrow pe$ in a shorthand notation). Because the energy distribution of electrons in β decay is not constant, a two-body decay is insufficient. Fermi adopted a hypothesis of Pauli that a third particle, of very low mass and uncharged, also participated in the interaction. He proposed a direct four-fermion interaction in which the fermions could be created or destroyed, motivated by photon emission in atomic transitions. His unconventional proposals proved correct; the neutrino was discovered in 1942 via the inverse process of β decay [12]. While his theory must be extended to describe higher energy interactions (and expanded to consider the parity violations, discussed in Section 1.4), it is a low-energy limit of the modern theory of the weak interaction. The connection of theories that are valid only for low-energy interactions to more complete theories has important implications for experimental inferences, which are increasingly challenges at higher energies. The techniques used in this thesis to study interactions beyond the directly accessible energy scale are heavily influenced by this work.

Such low-energy versions of a more complete theory are a valuable calculation tool, and give insight to the high-energy theory.

New equipment, such as the cloud chamber developed by Charles Wilson in 1911 [13], enabled extensive characterization of radioactive decays. In addition, it opened our eyes to the radiation bombarding the earth from outside the atmosphere. The discoveries of the positron in 1932 [14, 15] and the muon in 1936 [16] by Anderson

and others were made from such studies. Both had important implications for the burgeoning field of particle physics. The positron, a mirror image of the electron but with positive charge, had been predicted a year earlier by Dirac in his relativistic theory of the fermions [17]. His proposal of *antiparticle* partners to the observed matter particles, with equal mass but opposite quantum numbers, was too outlandish for many to take seriously but proved prescient. Further studies demonstrated that the muon behaves like a heavier partner to the electron, the first sign of *generations* of fermions. Because the particle seemingly played no role in the structure of the atomic, the existence of the muon was completely unexpected. Given the new breadth of the seemingly-fundamental particles observed, not all closely associated with the atomic theory, the field of particle physics was fully its own domain by the 1940s, with the purpose of characterizing the fundamental objects of the universe, their properties, and interactions.

1.2 Particle scattering experiments

In addition to studying the products of radioactive decays themselves, Earnest Rutherford realized the energetic decay products of radioactive nuclei could themselves be used as probes of other nuclei. By directing the helium nuclei (referred to as α particles) emitted from radioactive radium decay towards gold foil and measuring the deflection angles, he and his collaborators demonstrated that the deflection patterns of the scattered α particles were consistent with a concentrated positive charge at the center of atoms, the atomic nucleus [6]. Inferring information about the properties and interactions of particles based on their behaviour when scattered—referred to as a “scattering experiment”—is still central to particle physics experiments today.

This thesis concerns measurements of proton-proton (pp) scattering. A study of

scattering experiments in which protons colliding at sufficiently high energy to overcome the electromagnetic repelling force to allow a direct interaction of the quarks and gluons that compose the protons is performed. Fig. 1.1 gives a sketch of the scattering experiment approach to an important process in pp collisions, the details of which are not important: the production of the Higgs particle in decays to two photons (γ). Protons of known energy and momentum are brought to collisions and the outgoing particles—photons, in this example—are detected. Measurements of the outgoing particles (the photons), along with knowledge of the incoming particles, are used to infer information about the unseen interaction. As discussed in Chapters 2 and 4, such diagrams are not just useful visualizations of the interaction (in this case, the production and decay of the Higgs boson); they can be connected to mathematical expressions used to predict explicit properties of the outgoing particles. In this example, predictions of photon properties both with and without the existence of the Higgs boson can be compared to the experimental observation, which conclusively favors the with-Higgs-boson scenario. In collisions of interest to this work, the quark and gluon interactions produce massive vector bosons of the weak force that themselves interact. The consistency of these events with their expected properties is measured and used to test theories of new particles or new interactions.

Because high-energy particle interactions are driven by quantum mechanics, they are stochastic in nature. Therefore, a single collision cannot be uniquely interpreted in terms of a given interaction (or “process”). Rather, features of a measured process, such as the consistency or inconsistency with a hypothetical interaction or the presence of a new particle, must be inferred from collective results. To this end, a critical metric of scattering processes is the scattering cross section. For a particle beam scattered off a target, or equivalently, two beams of particles brought to collision, the cross section is the effective area of the beam deflected by scattering from the target

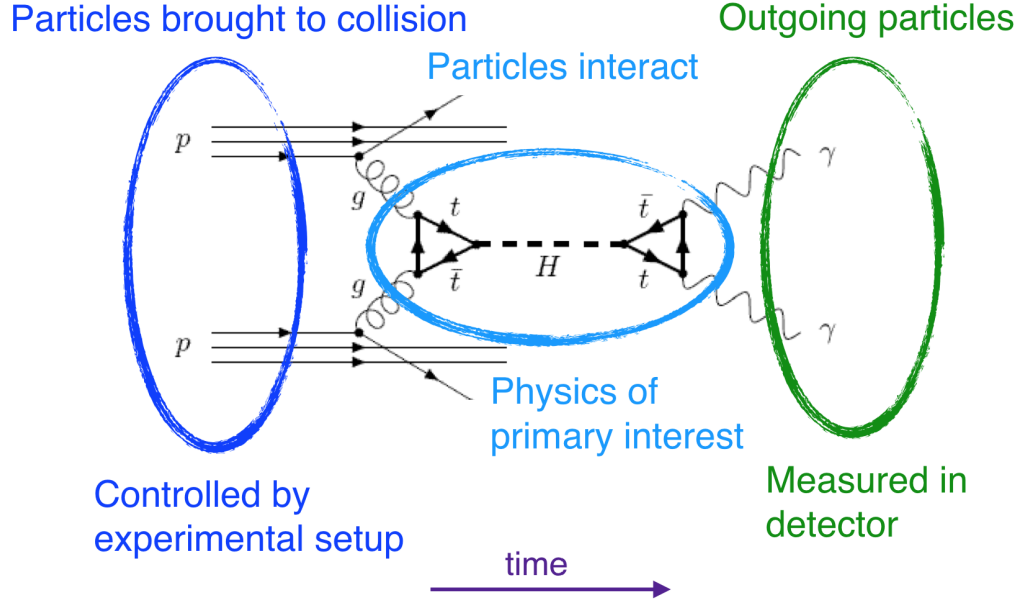


Figure 1.1: Illustration of Higgs boson production with decay to two photons in a pp scattering experiment.

(or other beam). For two beams with densities $\rho_A(x)$, $\rho_B(x)$ and length ℓ_A , ℓ_B , the cross section is defined as [18]

$$\sigma = \frac{\text{Number of scattering events}}{\ell_A \ell_B \int d^2x \rho_A(x) \rho_B(x)}. \quad (1.1)$$

For example, the cross section $\sigma_{pp \rightarrow e^+e^-}$ relates the number of events with two electrons produced in the scattering of two proton beams. The cross section is only unambiguously defined experimentally when it relates only the initial state and final states. However, it is common to discuss the cross section for, e.g., Higgs boson production, under the assumption that a theoretical extrapolation has been made in order to connect the measured particles (e.g., photons) to their production mechanism. Cross sections are expressed in units of area. It is customary to use the units of barns (b), where $1 \text{ b} = 10^{-24} \text{ cm}^2$.

It is also often useful to discuss not just the total cross section for all scattering to a final state, but to consider a *differential cross section*, dependent on properties of

the final state particles. For example, the differential cross section $d\sigma_{pp \rightarrow e^+e^-}/dE_{e^-}$ is measured by dividing the number of scattering events by properties of the outgoing electron energy E_{e^-} to produce a differential distribution of the production rate in terms of E_{e^-} . A measurement may also be restricted to consider only the scattering of outgoing particles into a specific geometrical region or above some threshold energy and momentum. Such a measurement is termed a *fiducial cross section*, restricted to a fiducial phase space defined by the geometrical and kinematic restrictions of the measurement.

1.3 Particle colliders and theoretical foundations

Rutherford continued to use the radioactive decays of elements as a means for particle scattering measurements through the early 1900s. Yet it was clear that a means of accelerating particles to collision at higher energies would be needed to further probe the structure of the nucleus. This requirement led to the invention of the first particle accelerators, which have come to define the field ever since. An electrostatic accelerator was developed at the Cavendish Laboratory in Cambridge, England by John Cockcroft and Ernest Walton and was used to split Lithium into Helium atoms in 1932 [19]. Concurrently in California, Ernest Lawrence developed the cyclotron to accelerate charged particles along a helical trajectory through electric and magnetic fields [20], which became a model for many future accelerator facilities. These inventions triggered an explosion in the field, marked by a continual effort to collide particles at higher and higher energies, capable of exploring the structure of matter and increasingly smaller distance scales.

Accelerators provided a new means to produce particles in the laboratory. Combined with improved tools for detection, they introduced a new era of exploration.

By the 1950s, particle physicists were faced with a “particle zoo” of seemingly fundamental objects spanning a huge range of masses. An underlying structure of the properties of many of these particles, now known as hadrons, was described by Gell-Man [21] and Zweig [22] in the early 1960s. They realized the properties of the many observed hadrons—including the proton and neutron—could be attributed to combinations of different elementary particles, which Gell-Man termed quarks. Whether these particles were physical or abstract entities was not immediately evident, but experiments at the Stanford Linear Accelerator a few years later showed that the proton has point-like constituents with the properties Gell-Man and Zweig predicted [23].

It remained unclear that the language of a quantum gauge theory, so successful in the formulation of the electromagnetic interaction, could also describe the strong or weak interactions. Both have crucial differences with respect to the electromagnetic interaction: weak interactions extend over a much smaller range, and the strong interactions has an inverted distance scaling, increasing at higher separation of the interacting objects while decreasing at very short ranges. A quantum gauge theory of the weak interactions was achieved in the 1960s with leading contributions from Glashow, Salam, and Weinberg [24]. Their theory unified the weak and electrodynamic interactions, with each communicated by the exchange of vector bosons. A crucial component of the electroweak theory is the presence of a scalar field that allows the vector bosons of the weak force, the $W^\pm$ and Z bosons, to acquire mass through electroweak symmetry breaking, further discussed in Chapter 2.

A description of the strong force as a quantum gauge theory (termed quantum chromodynamics, or QCD) followed, based on a more complex set of gauge fields. David Gross, David Politzer and Frank Wilczek overcame its computational complexity to demonstrate the required inverse distance scaling properties of the theory in 1973 [25, 26].

Many major discoveries that enabled and confirmed these theoretical breakthroughs are closely linked with technological advancements and collaborative efforts leading to more powerful accelerators, including

- The gluon, the vector boson of the QCD theory, discovered at the PETRA collider at Deutsches Elektronen-Synchrotron by the TASSO experiment [27], which provided conclusive support of QCD as a quantum gauge theory.
- The observations of the $W^\pm$ and Z bosons, the massive vector bosons of the weak interaction, discovered by the UA1 experiment at the SPS collider at the European Center for Nuclear Research (CERN) [28–31].
- The discovery of the top quark, the heaviest fundamental particle, by the D0 and CDF Collaborations at the Tevatron collider at Fermilab [32, 33].

The most powerful collider ever built, the Large Hadron Collider (LHC) at CERN, now allows the study of the most energetic particle collisions ever recorded in the laboratory, reaching concentrated energies only replicated in the most cataclysmic events in the universe. The LHC was planned, developed, and built over the course of decades, with personnel and funding from countries throughout the world. It enabled the discovery of the Higgs boson, the most recent fundamental discovery in particle physics. Its discovery in 2012 by the ATLAS [34] and CMS [35, 36] Collaborations is the last piece in the EW theory, confirmed almost 50 years after its original proposal by Higgs, Brout, Englert, Kibble, Hagen, and Guralnik [37–42].

The LHC accelerates proton beams to energies of 6.5×10^{12} terra-electronvolts (TeV), where $1 \text{ eV} = 10^{-12} \text{ TeV} = 1.602 \times 10^{-19} \text{ J}$, that are brought to collisions at a center-of-mass energy $\sqrt{s} = 13 \text{ TeV}$. In addition to accelerating protons to groundbreaking energy, the LHC delivers a very high rate of pp collisions. This allows very rare interactions to be studied, and is critical to obtaining statistically significant

results. In modern particle scattering experiments, the beam is not continuous, as suggested by Equation 1.1. It is helpful to express the properties of the beam particle density and effective length in terms of the luminosity $\mathcal{L}$ which captures the geometric properties of the beam such that

$$\text{Number of scattering events} = \sigma \mathcal{L}. \quad (1.2)$$

Luminosity has units of inverse area, customarily expressed in b^{-1} .

This work is based on studies of pp collisions delivered by the LHC and collected by the Compact Muon Solenoid detector in 2016, which are discussed in detail in Chapter 3 of this thesis. The precise definition of the luminosity for this experiment will also be discussed.

1.4 The particle content of the standard model

The experimental and theoretical work discussed here, as well as many other successful discoveries, measurements, and breakthroughs—combined with countless null results and unfulfilled hypotheses—have lead to the modestly-named standard model (SM) of particle physics. With the discovery of the Higgs boson, the SM is widely considered to be complete [43].

The particle content of the SM is depicted in Fig. 1.2, subdivided by the properties and quantum numbers of each particle. Matter in the SM is composed of fermions with half integer spin; all fermions in the SM have spin $1/2$. The spin 1 vector bosons, highlighted in red, mediate interactions between the fermions. Vector bosons are associated with the SM forces they mediate: the electroweak force consists of the electromagnetic force, mediated by the photon (γ), a neutral component of the weak force, mediated by the Z boson, a charged component of the weak force mediated by the W^+ and W^- bosons; the strong force is mediated by the gluon (g).

The particle masses, one of the most familiar properties, are also shown. The units of mass, GeV/c^2 , are motivated by the mass-energy relation $E = mc^2$; $1 \text{ GeV}/c^2 = 1.79 \times 10^{-27} \text{ kg}$. In this thesis, natural units (or God-given units, in the language of Ref. [18]) where $c = \hbar = 1$ are used. Here c is the speed of light, and $\hbar$ is the Planck constant. In this paradigm, the units of eV are used to express mass, energy, and momentum. Mass determines how particles interact with gravity, of which a quantum description has not yet been achieved. The description of gravity is outside the scope of the SM, and the gravitational forces on all particles produced at the LHC are many orders of magnitude lower than even the weak interaction and are not considered. The particle masses, which span orders of magnitude and do not follow a conclusive pattern, are not predicted by the SM but must be experimentally established. Predictive power with respect to particle masses would be a highly desirable feature of an extension to the SM.

Each particle in the SM has an associated antiparticle, with the same mass but opposite quantum numbers. Antiparticles are represented by explicitly specifying the charge, e.g. e^+ for the positron, or with an overbar, such as $\bar{q}$ for a generic antiquark. It is possible for a particle to be its own antiparticle, as is the case with the photon and Z boson, and possibly the neutrino [44]. Particles and antiparticles are otherwise expected to have equivalent properties, which has been experimentally validated, even for bound states such as antihydrogen [45]. The predominance of matter over antimatter in the universe is one of the most profound open questions in physics [46].

The fermions can be subdivided into “quarks,” highlighted in green in Fig. 1.2, and “leptons,” highlighted in purple, depending on whether they carry the charge associated with the strong force, known as “color.” There are three generations of fermions, divided by columns in Fig. 1.2. The first generation consists of regular

Standard Model of Elementary Particles

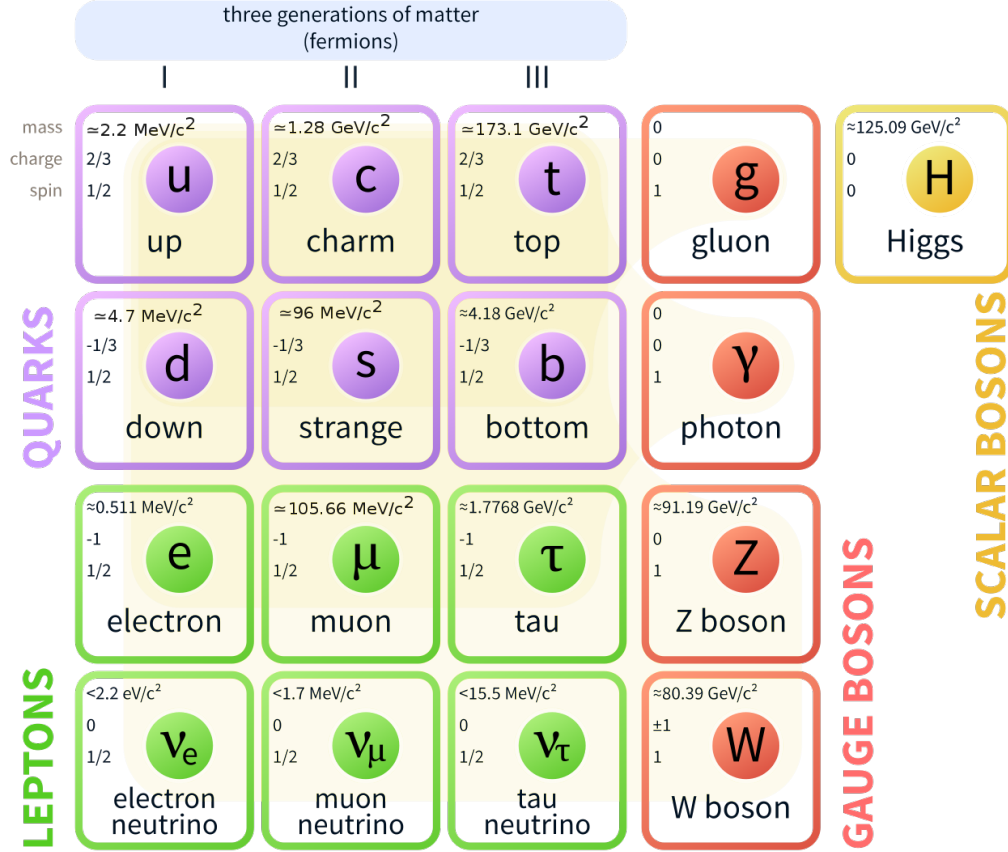


Figure 1.2: The experimentally established particle content of the universe.

matter: the up-quark (u), down quark (d), and electron (e^-). Second and third generation particles have similar properties but larger mass: the muon (μ) and tau (τ) leptons, and the charm (c) and strange (s) and top (t) and bottom (b) quarks. Lepton number, of which leptons (antileptons) carry $L = 1$ ($L = -1$) is conserved in the SM, as is baryon number—defined analogously in terms of the number of quarks (n_q) and antiquarks ($n_{\bar{q}}$) as $B = (n_q - n_{\bar{q}})/3$.

The vector bosons couple to the fermions that are charged under the force they mediate. All leptons but the neutrinos carry electric charge $-e$, where $e = 1.602 \times$

10^{-19} C. Therefore, they couple to the photon and participate in the electromagnetic force. Quarks carry electric charge $+2/3e$ or $-1/3e$. The $+2/3e$ quarks, shown in the top row of Fig. 1.2, are called up type, whereas the charge $-1/3e$ quarks are called down type (second row). Because the photon is massless, the range of the electromagnetic force is infinite. Its strength decreases with the separation of the interacting particles.

Unlike leptons, quarks carry color charge. Therefore, they couple to the gluon and participate in the strong force. Like the photon, the gluon is massless, yet unlike the electromagnetic force, the strength of the strong interaction increases with distance. Consequently, the quarks form bound states of short distance pairs or triplets of quarks, bound together by gluon exchange. There are three types of color charge: red, blue, and green (and respective anticolor, carried by antiparticles). Bound states of quarks are always “colorless,” either consisting of three quarks each of a different color, or a color-anticolor pair. We refer to $q\bar{q}$ bound states as mesons, the most familiar being the pions, formed from pairs of the u and d quarks. The $qq'q''$ states are known as baryons. They include the familiar proton, uud, and neutron, udd.

All SM fermions participate in the weak interaction. As its name suggests, the weak force is significant only over very short distance scales, a consequence of the fact that the W and Z bosons are massive. Two separate charges are associated with the weak force, weak hypercharge and weak isospin. Because the Z boson is uncharged, it communicates weak interactions of particles in which the quantum numbers are unchanged. Conversely, the $W^\pm$ bosons carry weak hypercharge and weak isospin. Interactions mediated by the $W^\pm$ bosons do not conserve quark flavor, that is, an interaction of quarks from one generation may give rise to quarks of another generation. The degree of quark mixing between generations is governed by the Cabibbo–Kobayashi–Maskawa (CKM) matrix, which must be determined experimentally. A

striking feature of the charged weak interaction is the violation of parity, or mirror symmetry, first observed experimentally by Wu [47] by measuring the β decays of cobalt-60 (based on a proposal by Lee and Yang [48]). As spin 1/2 particles, fermions have two possible possible *chirality* states, an abstract quantity that is equivalent to the particle helicity, or the relative alignment of the particle momentum and spin vectors, in the case of massless particles. Only left-handed fermions and right-handed antifermions participate in the interactions mediated by the W boson, which leads to the characteristic parity violation.

The Higgs particle, shown in yellow, has a unique role in the SM. It arises as a consequence of the spin-0 Higgs field, the only fundamental scalar in the SM. The presence of the Higgs field leads to the W and Z bosons obtaining mass, through a process known as electroweak symmetry breaking (EWSB). Fermions are also believed to acquire mass through interactions with the Higgs field, which has been experimentally verified for the heaviest quarks and leptons [49]. Understanding the nature of the Higgs bosons and the scalar sector of the SM is a major focus of particle physics today.

The underlying mathematical framework used to understand the quantum properties of these particles and their interactions, as well as an explicit discussion of EWSB, will be presented in Chapter 2.

1.5 Overview of this work

This thesis presents measurements of particle production in pp collisions at the CERN LHC. Collisions giving rise to a total of three electrons or muons, two forward “jets” of clustered hadronic particles, and an imbalance of transverse momentum—characteristic of undetected neutrinos—are selected to isolate contributions involving

the simultaneous production of a W^+ or W^- and a Z boson, the heavy particles that communicate the weak force. Selected collisions of interest, referred to as events, are used to measure the rate of production of processes predicted by the standard model (SM) of particle physics, the most complete theoretical expression of the known particles and forces of the universe, and to search for hypothetical extensions of the SM.

A dedicated search for a rare, previously unobserved SM process, referred to as electroweak (EW) WZ production (or EW WZ production), is presented. This class of processes includes direct interactions of the $W^\pm$ and Z bosons, which are precisely predicted by the SM. The presence and production rate of the EW WZ process is intimately connected to the mathematical structure of the weak interactions of the massive vector bosons, as well as the phenomenon of EW symmetry breaking (EWSB) [50].

The discovery of a scalar boson with couplings consistent with those of the SM Higgs boson by the ATLAS and CMS Collaborations [34–36] provides evidence that the W and Z bosons acquire mass through the Brout-Englert-Higgs mechanism [37–42]. However, current measurements of the Higgs boson couplings [51, 52] do not preclude the existence of scalar isospin doublets, triplets, or higher isospin representations alongside the single isospin doublet field responsible for breaking the EW symmetry in the SM [53, 54]. In addition to their couplings to the Higgs boson, the non-Abelian nature of the EW sector of the SM leads to quartic and triple self-interactions of the massive vector bosons. New interactions or particles not predicted by the SM (generically referred to as physics “beyond the SM,” or BSM physics) in the EW sector is expected to include interactions with the vector and Higgs bosons that modify their effective couplings. Characterizing the self-interactions of the vector bosons is thus of great importance.

This thesis presents the first study of EW WZ production performed by the CMS Collaboration [55]. Selected events are used to place constraints on BSM physics in terms of explicit models predicting charged Higgs bosons, decaying to a $W^\pm$ and a Z boson, and on generalized models of new interactions modifying in the rate of production, or kinematic distributions, of selected events.

Chapter 2

Diboson production and scattering

2.1 Mathematical formalism of the standard model

Guided by experimental tests and discoveries, we have built a description of all the known interactions of the SM particles in the language of gauge quantum field theory (QFT). An in-depth discussion of QFT and the mathematical expression of the SM is the subject of many textbooks, e.g., Refs. [18, 56–59]. An overview of the ideas most central to this thesis is presented here, following the discussion of Refs. [18, 50].

Particles in QFT arise as excitations in quantum fields. The equations of motion and interactions for the quantum fields of the SM can be extracted from the SM Lagrangian, which can be divided into terms governing the propagation of the free fields of the spin 1/2 fermions (quarks and leptons), the spin-0 Higgs field, and the vector boson fields:

$$\mathcal{L}_{SM} = \mathcal{L}_{\text{gauge}} + \mathcal{L}_{\text{leptons}} + \mathcal{L}_{\text{quarks}} + \mathcal{L}_{\text{scalar}} . \quad (2.1)$$

In this expression, $\mathcal{L}_{\text{gauge}}$ is a function of the spin 1 vector fields. The $\mathcal{L}_{\text{quark}}$ term

includes kinetic terms describing the propagation of free quarks as well as interaction terms coupling the quarks to the vector fields. The term $\mathcal{L}_{\text{leptons}}$ is built from similar kinetic terms and interaction terms coupling the leptons to the electroweak fields, reflecting the fact that the leptons do not participate in the strong interaction. For free fields fermions ψ , the kinetic terms of the Lagrangian are defined as

$$\mathcal{L} = \bar{\psi}(i\gamma^\mu\partial_\mu - m)\psi \quad (2.2)$$

where ψ is the spinor fermion field, γ^μ are the Dirac (or γ) matrices, and m is the mass of the fermion. The equations of motion derived from this expression give the well-known Dirac equation.

A fundamental property of the SM Lagrangian is its invariance under a class of transformations known as “gauge transformations.” The vector fields A_μ^a are referred to as *gauge fields* because of their role in ensuring that the Lagrangian remains invariant under a particular transformation, which may be respected by subsets of the fermionic fields. A gauge transformation is “generated” by the Lie group defined by a set of $n \times n$ matrices t^a . Infinitesimal unitary rotations of the fields of the form $\exp(iV(x))$, for an arbitrary $n \times n$ matrix $V(x)$ in the group, can be expressed in terms of the t^a matrices and an arbitrary function $\alpha(x)$ as $V(x) = 1 + i\alpha^a(x)t^a + \mathcal{O}(\alpha^2)$. Under this rotation, a gauge field transforms as

$$A_\mu^a \rightarrow A_\mu^a + \frac{1}{g}\partial_\mu^a + f^{abc}A_\mu^b\alpha^c. \quad (2.3)$$

Here g is the coupling constant of the theory and f^{abc} are the *structure constants* of the transformation, related to the commutator of the generating matrices as $[t^a, t^b] = if^{abc}t^c$. By introducing the *covariant derivative* D_μ associated with the gauge field A_μ ,

$$D_\mu = \partial_\mu - igA_\mu^a t^a. \quad (2.4)$$

and promoting the partial derivative of the Lagrangian in Equation 2.2, to a covariant derivative, the expression is naturally invariant under these transformations. Furthermore, the expression has further given rise to interactions between the fermion fields and the gauge fields. In this way, we conclude that the interactions of the fermions with the vector bosons arise as a consequence of the gauge symmetry of the theory.

The full SM Lagrangian respects transformations of the group $SU(3)_C \times SU(2)_L \times U(1)_Y$. In this expression, $SU(3)$ ($SU(2)$) is the Lie group of all unitary two-by-two (three-by-three) matrices with determinant 1, and $U(1)$ is the Lie group of complex numbers. The subscripts C , L , and Y indicate the fermion and gauge fields that are impacted by the rotation: the $SU(3)$ transformation concerns only objects with color charge (quarks and gluons), and the $SU(2)$ and $U(1)$ rotations concern weak isospin charge, carried only by the left-handed fermions, and weak hypercharge. The measurements in this thesis probe the structure of the electroweak force, so it is instructive to consider this in more detail.

The left-handed leptons form doublets under the $SU(2)$ isospin rotation:

$$L_\ell = \begin{pmatrix} \nu \\ \ell^- \end{pmatrix}_L. \quad (2.5)$$

A similar expression holds for the quarks:

$$L_q = \begin{pmatrix} u \\ d' \end{pmatrix}_L, \quad (2.6)$$

with corollary expressions for the charm and strange and top and bottom quarks. In this expression, d' indicates that the mass eigenstates of the down-type quarks d are not isospin eigenstates, rather, they are related by the 3×3 CKM mixing matrix. The left-handed isospin doublets carry isospin $I = 1/2$. All right-handed fermions carry isospin $I = 0$, that is, they are uncharged under the $SU(2)$ component of

the interaction, resulting in the characteristic parity violation of the weak force. The left-handed lepton doublets carry weak hypercharge $Y = -1$, whereas the left-handed quark doublets carry $Y = 1/3$. The right-handed leptons have $Y = -2$; right-handed up-type (down-type) quarks have $I = 4/3$ ($2/3$).

The EW theory consists of an two gauge fields: a field $\mathbf{B}_\mu$, which itself transforms as a vector under the $SU(2)_L$ symmetry, and an isoscalar field A_μ associated with the $U(1)_Y$ symmetry. Coupling constants g and g' corresponds to the isovector field $\mathbf{B}_\mu$ and isoscalar A_μ respectively. The generators of the $SU(2)$ symmetry are the Pauli isospin matrices, denoted τ , whereas the generator of $U(1)$ is 1. Combining Equation 2.3 and 2.2 gives the full expression for the term leptonic component of Equation 2.1,

$$\mathcal{L}_{\text{leptons}} = \bar{R}_\ell i\gamma^\mu \left(\partial_\mu + i\frac{g'}{2}A_\mu Y \right) R_\ell + \bar{L}_\ell i\gamma^\mu \left(\partial_\mu + i\frac{g'}{2}A_\mu Y + i\frac{g}{2}\tau \cdot \mathbf{B}_\mu \right) L_\ell. \quad (2.7)$$

A similar expression is appropriate for the quarks, with additional gauge fields (the gluons) enforcing the $SU(3)$ gauge invariance and leading to QCD interactions.

The free-field form of the EW gauge fields can be expressed in terms of the field strength tensors, defined as

$$F_{\mu\nu}^\ell = \partial_\nu B_\mu^\ell - \partial_\mu B_\nu^\ell + g\epsilon_{jkl}B_\mu^j B_\nu^k, \quad (2.8)$$

where $\ell = 1, 2, 3$ for the components of the isospin vector field, and

$$f_{\mu\nu} = \partial_\nu A_\mu - \partial_\mu A_\nu \quad (2.9)$$

for the isospin scalar field. Then

$$\mathcal{L}_{\text{gauge, EW}} = -\frac{1}{4} \sum_\ell F_{\mu\nu}^\ell F_\ell^{\mu\nu} - \frac{1}{4} \sum_\ell f_{\mu\nu} f^{\mu\nu}. \quad (2.10)$$

The field strength tensor of the B_μ fields mixes components of the isospin vector. This is because the associated group symmetry, $SU(2)$, is *non-Abelian*, that is, generators

of the theory do not commute. This non-Abelian nature leads to a more complex structure to the $SU(2)$ component of the interaction, including the self-interactions of the vector bosons associated with the gauge fields, which are probed in this thesis.

Note that in both Equation 2.7 and 2.10 we have neglected the mass term of the form $m|\psi|^2$. In both cases, such a term would not respect the gauge symmetries of the EW theory. The importance of the mass of the $W^\pm$ and Z bosons, and its relationship to the strength of the weak force, has already been emphasized. Furthermore, masses of the fermions have long been experimentally established. Clearly, this cannot be the correct theory of nature.

2.2 Electroweak symmetry breaking

The solution to this seemingly insurmountable issue with the theory comes from the $\mathcal{L}_{\text{scalar}}$ term of Equation 2.1 through the phenomena of EWSB, which provides a mechanism for the underlying symmetry of a fundamental theory to be hidden from its observable properties. Consider a complex isospin doublet scalar field ϕ , with $I = 1/2$ and $Y = 1$, with self interactions defined by the Lagrangian

$$\mathcal{L}_{\text{scalar}} = (D_\mu \phi)^\dagger (D^\mu \phi) + \mu^2 \phi^\dagger \phi - \lambda^2 (\phi^\dagger \phi)^2. \quad (2.11)$$

Here D_μ is the covariant derivative of the EW theory, which we infer from Equation 2.7 to be

$$D_\mu = \partial_\mu + i\frac{g'}{2}A_\mu Y + i\frac{g}{2}\boldsymbol{\tau} \cdot \mathbf{B}_\mu. \quad (2.12)$$

If the parameters μ and λ in the potential of Equation 2.17 are real, the minimum of the potential is not at 0, but at $v = \pm\mu/\lambda$, and the vacuum state of the field, or vacuum expectation value (vev) $v = \langle \phi | \phi \rangle_0$, is nonzero. The two components of the doublet field are complex, so ϕ has a total of four degrees of freedom. However, due

to the gauge symmetry, we are free to select a gauge in which only one component of the field must be explicitly given. If we expand this term around the vacuum expectation value, then

$$\phi = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + h(x) \end{pmatrix}. \quad (2.13)$$

Inserting this new form of the field into Equation 2.17 has a dramatic effect: coupling terms between the h field and vector boson fields are present, as are new terms quadratic in the vector boson fields. Apparently, the explicit choice of the vev has broken the gauge symmetry and introduced mass terms for the vector bosons. The mass eigenstates are combinations of the $\mathbf{B}_\mu$ and A_μ fields:

$$\begin{aligned} W_\mu^\pm &= \frac{1}{\sqrt{2}} (B_\mu^1 \mp B_\mu^2) \\ Z_\mu &= \frac{gB_\mu^3 - g'A_\mu}{\sqrt{g^2 + g'^2}} = B_\mu^3 \cos \theta_W - A_\mu \sin \theta_W \\ A_\mu &= \frac{g'B_\mu^3 + gA_\mu}{\sqrt{g^2 + g'^2}} = B_\mu^3 \sin \theta_W + A_\mu \cos \theta_W. \end{aligned} \quad (2.14)$$

where the Weinberg angle θ_W has been introduced, defined as $g' = g \tan \theta_W$. Substituting these redefinitions of the vector and scalar fields of Equation 2.14 and 2.13, we see terms of the form

$$\mathcal{L}_{\text{gauge,m}} = -\frac{v^2 g^2}{4} W_\mu^+ W^{-\mu} - \frac{v^2 (g^2 + g'^2)}{8} Z_\mu Z^\mu. \quad (2.15)$$

The new massive vector states are exactly those of the $W^\pm$ and Z bosons, with masses $m_W = gv/2$ and $m_Z = v\sqrt{g^2 + g'^2}/2$, experimentally established to be $m_W = 80.385 \pm 0.015 \text{ GeV}$ and $m_Z = 91.1876 \pm 0.0021 \text{ GeV}$ [49]. The component of the scalar not absorbed by the gauge freedom is the Higgs field. It gives rise to the Higgs boson, with mass $m_H = \sqrt{2}\mu$, experimentally observed and established as $m_H = 125.09 \pm 0.24 \text{ GeV}$. The A_μ field remains massless, as required by the photon.

Its coupling constant is $e = gg'/\sqrt{g^2 + g'^2}$, with the conserved quantum number charge, $Q = I^3 + Y/2$, where I^3 is the third component of weak isospin.

In terms of the massive vector fields, the Lagrangian contains self-interactions of the vector fields and the Higgs fields. The couplings of these fields are exactly specified in terms of the coupling constants of the $SU(2) \times U(1)$ fields and the vector boson masses. Specifically, the VH and W WV couplings are

$$\begin{aligned} \mathcal{L}_{WWVV} = & -\frac{g^2}{4} \left\{ [2W_\mu^+ W^{-\mu} + (A_\mu \sin \theta_W - Z_\mu \cos \theta_W)^2]^2 \right. \\ & - [W_\mu^+ W_\nu^- + W_\nu^+ W_\mu^- \\ & \left. + (A_\mu \sin \theta_W - Z_\mu \cos \theta_W)(A_\nu \sin \theta_W - Z_\nu \cos \theta_W)]^2 \right\}, \end{aligned} \quad (2.16)$$

$$\mathcal{L}_{HV} = \left(gm_W H + \frac{g^2}{4} H^2 \right) \left(W_\mu^+ W^{-\mu} + \frac{Z_\mu Z^\mu}{2 \cos^2 \theta_W} \right). \quad (2.17)$$

It is important to note that the quartic interactions, given in Equation 2.16, are proportional to the square of the electroweak coupling g . As described in the following sections, g is significantly less than 1, and they are naturally suppressed with respect to many couplings electroweak processes.

Lastly, we note that Yukawa interaction terms between the scalar and fermion fields are gauge invariant:

$$\mathcal{L}_{\text{Yukawa}} = -\zeta_\ell [(\bar{L}_\ell \phi) R_\ell + \bar{R}_\ell (\phi^\dagger) L_\ell]. \quad (2.18)$$

Inserting the scalar field of Equation 2.13 into this expression gives rise to quadratic mass terms without violating the symmetry of the fundamental Lagrangian. Thus the Higgs field can also give rise to Fermion masses in gauge invariant way. Measurements of Higgs boson decays to fermions have provided strong evidence that this is indeed the mechanism by which the heaviest leptons obtain mass.

2.3 Particle scattering and perturbative calculations

With the full SM Lagrangian in hand, the principle of least action can be used to derive the equations of motion of the SM fields. For the kinetic terms, exact solutions can be obtained, however, closed-form equations of motion do not exist when the interaction terms are also relevant. The strength of these interaction is set by the coupling constants g_i associated with the gauge fields, as described for the electroweak interaction in Section 2.1. If the couplings, and therefore the interactions, can be considered small, they can be treated as a perturbation of the free-field propagation. Because field interactions are local, associated with a range of the interaction communicated by the gauge field, this treatment is particularly well suited for scattering experiments. Free, non-interacting particles of the fields are brought together in collision from a large separation, where the distance scale of the interaction is inversely proportional to the energy of the collision. After interaction, free particles, possibly of different type than those brought to collision, emanate from the interaction point.

The allowed interactions in the SM are illustrated in Fig. 2.1. Lines represent propagating particles, and vertices represent interactions, arising from the interaction terms in the SM Lagrangian. These illustrations, known as Feynman diagrams, can be pieced together to give a diagrammatic picture of a scattering interaction. Fig. 2.2 gives two illustrative qq interactions—a component of the pp collision at the LHC—mediated by the W and Z bosons, leading to free muons, an electron, and an anti-electron neutrino.

Feynman diagrams are not just useful as an illustration, they represent terms in the perturbative expansion used to model interactions, specifically, elements of the

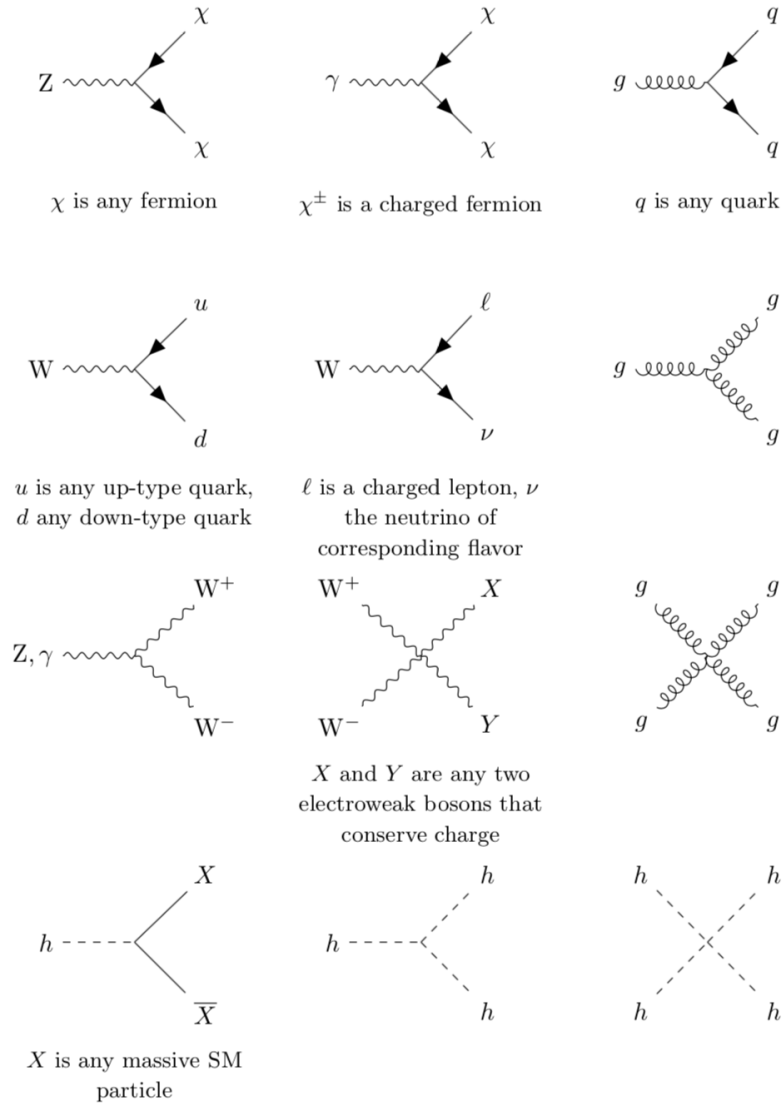


Figure 2.1: Interactions allowed in the SM, excluding those involving the Higgs field. Reproduced from Ref. [60].

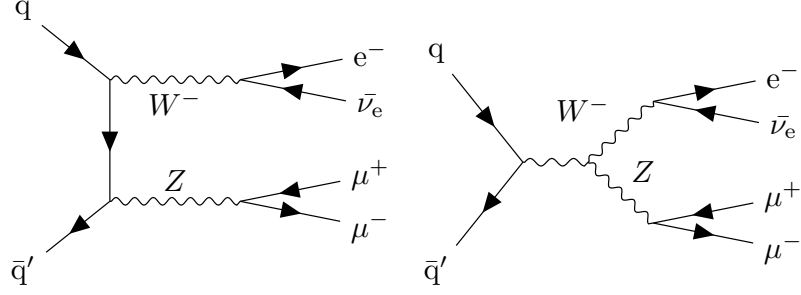


Figure 2.2: Feynman diagrams illustrating the $q\bar{q} \rightarrow e\bar{\nu}_e\mu^+\mu^-$ process, proceeding via quark and lepton couplings to the W and Z bosons.

scattering matrix that connects the initial and final state in a scattering experiment. So-called Feynman rules are used to associate the fields and interactions of a diagram with the mathematical formalism, giving mathematical expressions for the matrix elements connecting the outgoing particle four momenta and quantum numbers to the incoming particle properties. Observables, including the total rate of production for a final state given the initial state, are dependent on the square of the scattering matrix elements.

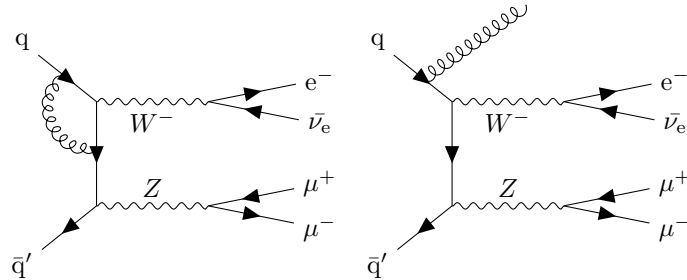


Figure 2.3: Feynman diagrams illustrating the $q\bar{q} \rightarrow e\bar{\nu}_e\mu^+\mu^-$ process with higher-order QCD couplings.

The diagrams in Fig. 2.2 are not the only possible set of interactions connecting the $q\bar{q}$ initial state to the $e\bar{\nu}_e\mu^+\mu^-$ state. More intricate arrangements of the quark and gluon lines are possible, as shown in Fig. 2.3 (left). Whereas the diagrams in Fig. 2.2 have four vertices of fields coupling via the EW interaction, this diagram has two additional QCD couplings of the gluon and quarks. Each vertex contributes and

additional factor of the EW or QCD couplings in the mathematical expression of the scattering matrix element. If the coupling term is $\ll 1$, the dominant contribution will come from diagrams such as those shown in Fig. 2.2, and it may be possible to neglect the contribution of Fig. 2.3 (left). Perturbative calculations rely on this approximation by restricting a calculation to a maximum “order” $\mathcal{O}(\alpha^n)$ in the coupling $\alpha \propto g^2$, by only considering contributions to a process that have at most a dependence α^n .

While the diagram in Fig. 2.3 (right) does lead to a $e^-\bar{\nu}_e\mu^+\mu^-$ state, it also produces a final-state gluon. It is tempting to consider this contribution as a separate process, however, due to the nature of the strong force, such an approach is not favorable experimentally. Free gluons are not observable, rather, they hadronize into baryons and mesons, which do not correspond trivially to the radiated partons. In addition, for gluons radiated at a small angle or low momenta, it is impossible to experimentally distinguish whether two or one quarks or gluons are present. For this reason, many leptonic measurements at the LHC are *inclusive* in the number of associated hadronic particles. That is, we consider $pp \rightarrow WZ$ production to include all states producing the $e\bar{\nu}_e\mu^+\mu^-$ with any number of associated hadronic objects.

Depending on the process considered and the accuracy needed for a calculation, neglecting the terms of “higher order” may not be appropriate. While the Feynman rules are equally applicable for these terms, additional contributions arise due to the presence of closed loops, as seen for the gluon and quark lines in Fig. 2.2 (right). These individual diagrams give results that predict infinite rates, therefore, they cannot correspond to experimental observations. Fortunately, the situation can be recovered when combining with other diagrams contributing to the higher-order calculation. A finite contribution can be isolated and used to describe experimental measurements with incredible accuracy. Infinite terms remain, but they can be associated with experimentally-established terms—in particular, the masses and cou-

plings of the theory. This procedure, known as *renormalization*, is closely connected to the energy scale of the process and the structure of the fields themselves. In this approach, the coupling constant captures the neglected contributions from diagrams contributing at higher perturbative orders, becoming an effective coupling at the energy scale considered. The effective coupling is a function only of the energy scale probed. Furthermore, the energy dependence of the effective coupling is a fundamental aspect of the theory that can be calculated to high accuracy. Thus, once the effective coupling is experimentally established for some assessible energy scale in a well-measured process, it can be applied in many other calculations.

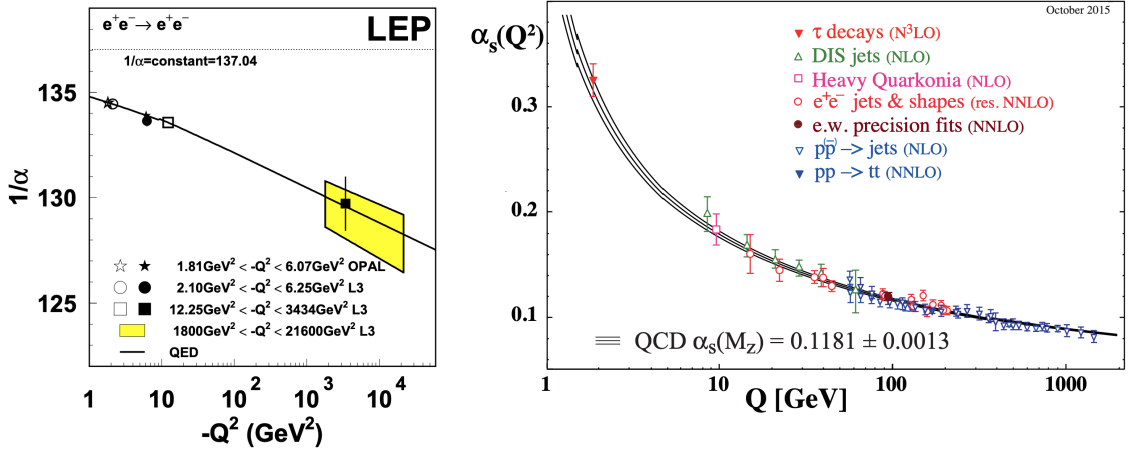


Figure 2.4: The energy scale dependence of the effective EW (left, reproduced from Ref. [61]) and QCD (right, reproduced from Ref. [49]) couplings. The theoretically predicted bands and discrete points established experimentally are shown.

The energy dependence of the effective coupling constants for the EW and QCD theories are shown in Fig. 2.4. As shown, the nature of the SU(3) and SU(2) forces leads to a striking difference in energy dependence between the two theories: the EW coupling strength increases with energy, whereas the QCD coupling decreases, leading to the confinement of quarks and gluons into bound states, such as the proton and neutron, at low energies. At the energy scale probed in LHC collisions, the QCD

coupling is sufficiently small to enable perturbative calculations, however, higher-order corrections are often significant: in diboson processes, typical corrections are $\mathcal{O}(100\%)$ from the lowest order in perturbation theory (leading order, or LO) to the next-to-leading (NLO) order, and 5–10% from NLO to next-to-next-to-leading order (NNLO) [62]. The EW coupling is relatively small compared to the QCD coupling at energies accessible to colliders, so higher-order EW corrections are often not essential to accurate calculations. Perturbative calculations are discussed in more detail in Chapter 4.

2.4 The WZ production at the LHC

At low energy, the proton can be viewed as a bound state of the quarks uud. However, at higher energy, the content of the proton appears more complex. High-energy quarks can radiate gluons, which can split to $q\bar{q}$ pairs. In the high-energy pp collisions at the LHC, the proton is a source of gluons, and all flavors of quarks but the heaviest, the top quark. Therefore, the initial state of fundamental particles in the interactions studied at the LHC is always composed of quarks and gluons, collectively dubbed “partons.” The energy of the colliding protons is controlled, but the flavor and energy of the interacting partons can only be stochastically determined based on these properties, as discussed in Chapter 4.

The dominant modes of $pp \rightarrow WZ \rightarrow \ell'\ell'\ell\nu$ (where $\ell, \ell' = e, \mu$) production are shown in Fig. 2.2. Equivalent diagrams are possible with the W and Z bosons coupling to quarks, or the Z boson to neutrinos. Many aspects of the processes are similar, and it is often justified to treat the $pp \rightarrow WZ$ and WZ decay to leptons, neutrinos, or quarks separately. For this reason, we refer to “WZ production,” keeping in mind that the W and Z bosons have an extremely short lifetime and are only observable via their

decay products. The results in this thesis exploit the leptonic decay channels, $W^- Z \rightarrow e^+ e^- \mu^- \bar{\nu}_\mu$, $W^+ Z \rightarrow e^+ e^- \mu^+ \nu_\mu$, $W^- Z \rightarrow \mu^+ \mu^- e^- \nu_\mu$, and $W^- Z \rightarrow \mu^+ \mu^- e^+ \nu_e$.

Studies of WZ production at the LHC provide an important test of the SM. As shown in Fig. 2.2 (right), the process is sensitive to the charged WZZ coupling, which arises due to the non-Abelian nature of the electroweak theory, and is exactly predicted in the SM [63]. Additional charged resonances, or other modifications of the EW sector of the SM, would modify this process and adjust the rate. Moreover, the process is known to be sensitive to higher-order corrections [64]. Measuring this process with sufficient accuracy to test state-of-the-art calculations, and to determine consistency with the SM or uncover hints of physics not predicted by the SM—which we refer to generically as “new physics”—is of great interest to the LHC experimental program.

As mentioned in the previous section, it is often advantageous to consider measurements inclusive in the number of hadronic particles. It is possible, however, to make exclusive measurements, provided that we are careful in how we associate the number of hadronic objects to the number of partons in a Feynman diagram. The solution is to cluster the partons, or hadrons, using algorithms that are well-behaved for low energy partons. We refer to these clustered hadronic objects as “jets,” (j) and a close correspondence between experimental measurements and theoretical predictions has been demonstrated for several so-called jet-clustering algorithms [65]. The specific algorithms used in this analysis will be discussed in Chapter 5.

The focus of this thesis is the process $pp \rightarrow WZjj$, an important subcomponent of the inclusive $pp \rightarrow WZ$ process. While it is often useful to relate the $WZ + 2$ partons calculation and the $WZ + 2$ jets ($WZjj$) experimental state, it is important to note that the relationship is most appropriate when using a common jet clustering algorithm. The dominant contribution to the $WZjj$ state is a higher-order correction

to the diagrams of Fig. 2.2, where gluons are radiated from the incoming quarks. Because of the additional QCD vertices, these diagrams contribute to the WZjj state at $\mathcal{O}(\alpha_s^2\alpha^2)$, compared with $\mathcal{O}(\alpha_s^0\alpha^2)$ for the dominant contributions to the $pp \rightarrow WZ$ process. At this order, contributions from qg and gg initial states also contribute, slightly enhancing the total WZjj production rate.

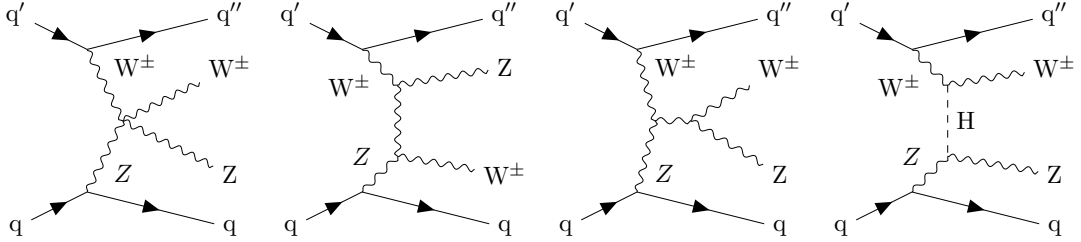


Figure 2.5: Representative Feynman diagrams for EW WZ production in the SM, including those with quartic WWZZ interactions (left) and via scattering with the Higgs boson (right).

In addition to these contributions to the WZjj production, another contribution at $\mathcal{O}(\alpha_s^0)$ contributes, of order $\mathcal{O}(\alpha_s^4)$. This subclass of processes with only EW couplings at tree level has a rich and important phenomenology. It includes contributions from vector boson scattering (VBS), where vector bosons are radiated from the incoming quarks before interacting, as illustrated in Fig. 2.5. The VBS processes form a distinct experimental signature characterized by the W and Z bosons with two forward, high-momentum jets, arising from the hadronization of two quarks. The previously discussed contributions to the WZjj state that proceed via QCD radiation of partons from are referred to as QCD-induced WZjj production (or QCD WZ).

2.5 Probing the electroweak sector through vector boson scattering

As illustrated in Fig. 2.5 (right), EW WZ production is sensitive to the quartic couplings of the vector bosons, as well as the interactions of the vector bosons and the Higgs boson. Additional interactions in the EW sector, such as heavy new scalar Higgs bosons, would lead to additional contributions, modifying the effective rate of production of EW WZ events.

The existence of the Higgs boson has strong implications for VBS VV production. In perturbative calculations of longitudinally polarized VV scattering, the production rate involving the diagrams of Fig. 2.5 that do not involve the Higgs couplings increase with the scattering energy without bound. In 1987, an analysis of VBS production of $W+W^-$ in pp collisions by Lee, Quigg, and Thacker demonstrated that the presence of a scalar boson is necessary to regulate this component to a finite rate [66]. By analyzing the partial wave amplitudes of the process, they derived a unitary bound on for the process that stipulated $m_H \lesssim 1$ TeV, providing a convincing argument for the Higgs boson to be discovered at the future LHC. While the observed SM Higgs is sufficient to regulate the process, further modifications could disturb this cancellation, leading to large deviations from the SM predictions.

This thesis presents a measurement of the EW WZ process under the assumption that its rate and characteristics are those of the SM prediction. We also search for deviations from the SM in both the rate of production of events characteristic of the EW WZ process, and in differential distributions sensitive to the scattering energy of the interaction and the mass of new scalar particles. We interpret the results in terms of both explicit models predicting charged Higgs bosons, shown in Fig. 2.6 (left), and in the generalized framework of dimension-8 effective field theory, illustrated in

Fig. 2.6 (right). In this diagram a set of unknown interactions are represented, which result in a modification of the effective quartic vertices. The charged Higgs and EFT models we investigate are complimentary: the former is applicable in a regime where the collision energy allows the resonant Higgs to be directly produced, whereas the latter is meaningful when the scale of new physics exceeds the LHC collision energy. The models we consider are described in more detail in the following sections.

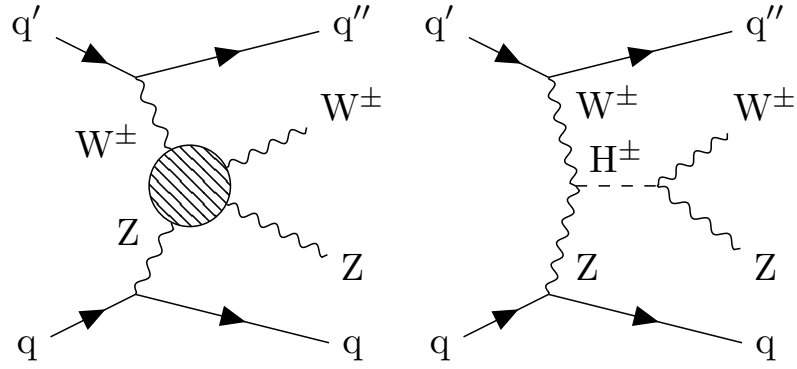


Figure 2.6: Feynman diagrams illustrating hypothetical new physics in the EW WZ process. Charged Higgs boson production (left) is predicted in many BSM extensions. The aQGC approach (right) can be used to parameterize a large class of new interactions modifying the WWZZ process.

2.5.1 Additional Higgs bosons and the Georgi-Machacek model

In the SM, EWSB is realized by a single isospin doublet scalar field, as discussed in Section 2.2. This field is sufficient to give mass to the W and Z bosons, while resulting in a scalar H boson consistent with the observed boson with $m_H = 125 \text{ GeV}$. Given the complexity of the fermion sector of the SM, it is natural to ask if the universe contains a more complex structure of scalar fields than the single doublet. However, SM extensions by arbitrary SU(2) scalar fields are highly constrained by the parameter $\rho \equiv m_W^2/m_Z^2 \cos^2 \theta_W$, experimentally established to be very close to 1 [49].

For a theory with n scalar multiplets ϕ_i , with weak isospin I_i , weak hypercharge Y_i , and vacuum expectation value v_i , ρ is defined in terms of the quantum numbers of the fields by the following expression at tree level [67]

$$\rho = \sum_{i=1}^n \frac{I_i(I_i + 1) - \frac{1}{4}Y_i^2}{\frac{1}{2}Y_i^2 v_i}, \quad (2.19)$$

Extending the scalar sector of the SM by two rather than one doublet satisfies this expression. This model, known as the Two-Higgs-doublet model (2HDM), has been extensively studied, in part because the scalar sector of the well-known theory of supersymmetry is required to have at least two Higgs doublet fields. The 2HDM gives rise to a heavy Higgs boson, a pseudoscalar Higgs boson, and charged Higgs bosons $H^\pm$ in addition to a light SM-like Higgs boson. In this model, the $H^\pm$ bosons couple strongly to the leptons, and decays to bosons are strongly suppressed in the mass range probed by the LHC [68].

Alternatively, if the Higgs sector of the SM is extended by scalar SU(2) triplets, charged Higgs bosons with tree-level couplings to the W and Z bosons are generated by EWSB. Higgs triplets appear in left-right symmetric [69, 70], little Higgs [71–73], and some supersymmetric models [74, 75]. A simple extension of the SM with Higgs triplets that satisfies experimental constraints on ρ was proposed by Georgi and Machacek [76]. In the Georgi–Machacek (GM) model, the Higgs sector is extended by a real triplet with $Y = 0$ and a complex triplet with $Y = 2$ in addition to the usual SM doublet ($Y = 1$). After EWSB, this leads to a rich Higgs sector with intriguing phenomenology. The Higgs sector of the GM has ten physical states that are organized based on their transformation properties under the approximate global SU(2) symmetry of the scalar fields remaining after EWSB, referred to as custodial symmetry [77]. There are two singlets, one of which is associated with the observed Higgs boson, a triplet, and a fiveplet (H_5^{++} , H_5^+ , H_5^0 , H_5^- , H_5^{--}) that

is fermiophobic. If the mass of the triplet states is greater than the mass of the fiveplet, the only production mechanism for $H^\pm$ at the LHC is vector boson fusion (VBF, shown in Fig. 2.6, left). The production cross section is proportional to the variable $s_H^2 \equiv \sin^2 \theta_H$, where $\sin \theta_H = 2\sqrt{2}v_\chi/v$, the ratio of the vacuum expectation value (vev) of the triplet fields to the SM vev. The WZjj state probed in this thesis is well-suited for further probes of this model, which is significantly less constrained than models predicting charged Higgs bosons with predominantly leptonic decays.

2.5.2 Generalizing new interactions with effective field theory

Explicit models of charged Higgs boson production, such as those discussed in the previous section, provide a well-motivated set of characteristics to exploit in searches for deviations from the SM. In such a paradigm, expectations about the underlying structure of the theory allow explicit relationships between independent measurements to be drawn. However, when results favor the SM to a proposed extension, it may be useful to cast a wider net, in search of more subtle or more exotic extensions. Furthermore, if the mass scale of new physics—such as a heavy charged Higgs bosons—exceeds the range directly probed by 13 TeV pp collisions, a broad characterization of a deviation, or lack thereof, may be more meaningful than sensitivity to a particular model.

Effective field theory (EFT) is well suited for a broad characterization of new physics at a high energy scale. In EFT, the SM is assumed to describe the field content of the universe below some energy scale Λ . Additional fields may exist, but they are only directly accessible (e.g., can be produced at a collider) above Λ . Below this scale, their effect on the low-energy world can be parameterized in terms of the

known SM field content. The effective interactions are captured by extending the SM Lagrangian with additional terms of mass dimensions $d > 4$ that are built from the SM fields. The effective terms are suppressed by dimensionful couplings proportional to the $1/\Lambda^{4-d}$, leading to the expected energy dependence of the interactions. The treatment breaks down as the energy scale approaches Λ , but for energies $\ll \Lambda$ all effective interactions can be described.

The bottom-up approach to EFT described here, in which new interactions are parameterized in terms of the low-energy field content, is useful because it is consistent with a top down approach, where the low-energy behavior is extracted from a known full theory [78]. The most well-known example of this is the Fermi theory of the weak decay $n \rightarrow \bar{p}e^-\bar{\nu}_e$. A low energy description, in which the leptons, quarks, and neutrinos couple directly in a four-fermion interaction was originally proposed by Fermi without knowledge of the underlying theory. In the full EW theory, the interaction between a d quark of the neutron and an electron is mediated by the W^- boson, through dW^-u and $e^-W^-\bar{\nu}_e$ couplings of the EW theory. Because the momentum exchange p of the interaction is on the MeV scale, whereas $m_W = 80.4 \text{ GeV}$, the full expression for the interaction can be expanded in terms of the ratio p/M , to derive an approximate theory valid at low energy with the effective coupling $G_F = \sqrt{s}g^2/8m_W^2$ of Fermi's theory. In this sense, the low-energy EFT gives insight gives useful predictions within its energy regime and provides insight into the nature of the full theory.

In this analysis, we consider a set of dimension-8 EFT operators sensitive to the quartic interactions of the W , Z , and γ . Because an alternative, largely historical, approach to parametrizing these interactions involves directly modifying the coupling terms in the Lagrangian, we refer to this as a search for anomalous quartic gauge couplings (aQGC). Similarly, anomalous triple gauge couplings (aTGC), probed in

inclusive WZ production, can be parameterized by dimension-6 operators.

The set of operators considered in this analysis are shown in Equation 2.20.

$$\begin{aligned}
\mathcal{L}_{S0} &= \frac{f_{S0}}{\Lambda^4} [(D_\mu \Phi)^\dagger (D_\nu \Phi)] \times [(D^\mu \Phi)^\dagger (D^\nu \Phi)] \\
\mathcal{L}_{S1} &= \frac{f_{S1}}{\Lambda^4} [(D_\mu \Phi)^\dagger (D^\mu \Phi)] \times [(D_\nu \Phi)^\dagger (D^\nu \Phi)] \\
\mathcal{L}_{T0} &= \frac{f_{T0}}{\Lambda^4} \text{Tr} [\hat{W}_{\mu\nu} \hat{W}^{\mu\nu}] \times \text{Tr} [\hat{W}_{\alpha\beta} \hat{W}^{\alpha\beta}] \\
\mathcal{L}_{T1} &= \frac{f_{T1}}{\Lambda^4} \text{Tr} [\hat{W}_{\alpha\nu} \hat{W}^{\mu\beta}] \times \text{Tr} [\hat{W}_{\mu\beta} \hat{W}^{\alpha\nu}] \\
\mathcal{L}_{T2} &= \frac{f_{T2}}{\Lambda^4} \text{Tr} [\hat{W}_{\alpha\mu} \hat{W}^{\mu\beta}] \times \text{Tr} [\hat{W}_{\beta\nu} \hat{W}^{\nu\alpha}] \\
\mathcal{L}_{M0} &= \frac{f_{M1}}{\Lambda^4} [(D_\mu \Phi)^\dagger (D_\nu \Phi)] \times [(D^\mu \Phi)^\dagger (D^\nu \Phi)] \\
\mathcal{L}_{M1} &= \frac{f_{M1}}{\Lambda^4} [(D_\mu \Phi)^\dagger (D^\mu \Phi)] \times [(D_\nu \Phi)^\dagger (D^\nu \Phi)]
\end{aligned} \tag{2.20}$$

Here

$$\hat{W}_{\mu\nu} = \sum_j W_{\mu\nu}^j \frac{\sigma^j}{2}, \tag{2.21}$$

where $W_{\mu\nu}$ are the $\text{SU}(2)_L$ fields before EWSB and the values f_{Si} , f_{Ti} , and f_{Mi} are dimensionless couplings to the dimension-8 operators.

The operators probed are a subset of those outlined in Ref. [79], which proposes nine independent dimension-eight operators that assume the $\text{SU}(2) \times \text{U}(1)$ symmetry of the EW gauge sector as well as the presence of an SM Higgs boson. All operators are charge conjugation and parity conserving. The $\mathcal{L}_{Si}$ operators are a function of the scalar Higgs field. The $\mathcal{L}_{Ti}$ operators are functions of the gauge fields, and the $\mathcal{L}_{Mi}$ operators are built from a mix of the scalar and gauge fields. The quartic gauge coupling can be expressed in terms of these operators as

$$\mathcal{L}_{quartic} = f_{S0} \mathcal{L}_{S0} + f_{S1} \mathcal{L}_{S1}. \tag{2.22}$$

With $f_{S0} = -f_{S1} = f$, non-zero values of f give a rescaling of the SM coupling by $(1 + fv^4/8)$.

2.6 Previous results

2.6.1 Measurements of WZ production

The WZ production was first observed in $p\bar{p}$ collisions by the CDF [80, 81] and D0 [82] Collaborations at the Tevatron in 2006. The inclusive WZ production cross section in pp collisions has been measured in the leptonic decay modes by the ATLAS and CMS Collaborations at 7, 8, and 13 TeV [83–88]. A summary of these measurements is shown in Fig. 2.7. The summarized experimental results are obtained by measuring $pp \rightarrow WZ \rightarrow \ell'\ell'\ell\nu$ process, corrected for the independently-measured leptonic branching ratios (the percentage of events with $W \rightarrow \ell\nu$ and $Z \rightarrow \ell'\ell'$ to all other decay channels), to obtain a total rate for the process $pp \rightarrow WZ$. Theoretical predictions are shown at NLO in QCD, in which all contributions to the process up to $\alpha^2\alpha_s$ are considered in the calculation, and at NNLO, which considers contributions up to $\alpha^2\alpha_s^2$. The predictions are obtained using the programs MCFM [89, 90] and MATRIX [62, 64]. The data exhibit a slight preference for the higher-order prediction, demonstrating the importance of both high-precision measurements and calculations.

2.6.2 Measurements of WZjj production and electroweak

WZ production

Fig. 2.8 shows the distribution of the number of jets associated with the selected $\ell\ell'\nu$ state in data collected by the CMS experiment in 2015. This analysis, the first measurement of the WZ process at 13 TeV, was a critical first step towards the WZjj studies presented in this thesis. The expected contribution from the WZ process, along with the estimated backgrounds, are seen to be in good agreement with the observed data. This studies of this thesis are performed with a higher-luminosity

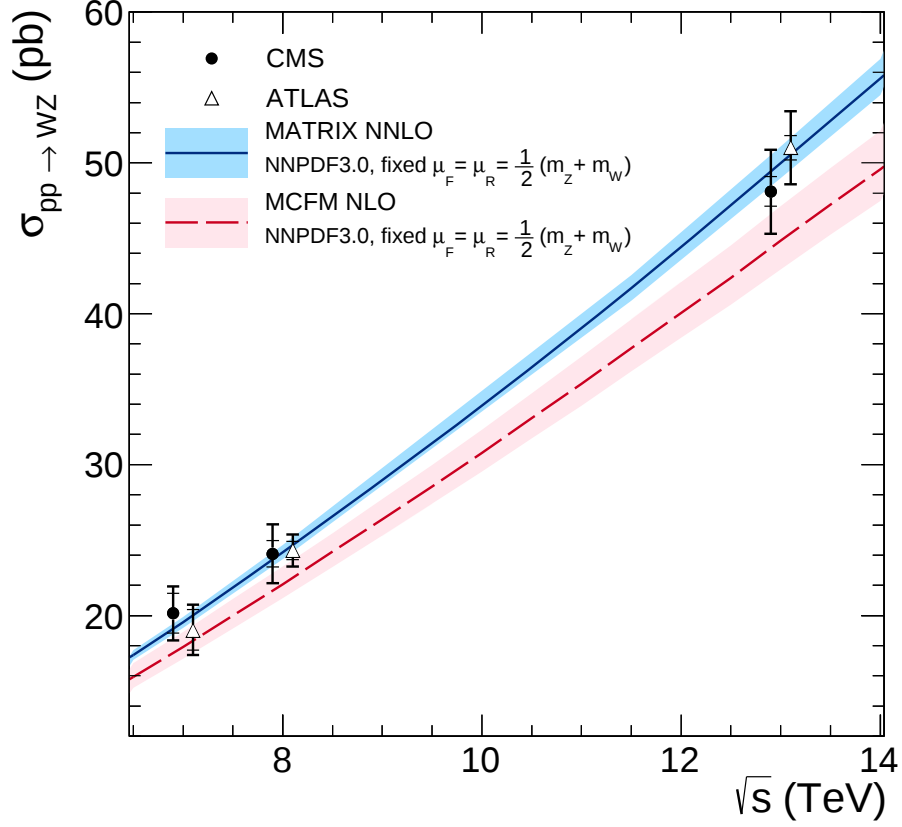


Figure 2.7: The measured $pp \rightarrow WZ$ production cross sections at 7, 8, and 13 TeV compared with theoretical predictions at NLO and NNLO in QCD.

data set, in order to exploit the kinematics variables of the selected jets to isolate the EWWZ and new physics contributions.

The first studies of VBS production of pairs of massive vector bosons—termed diboson production, or VV production, where $V = W, Z$ —were performed using in the $pp \rightarrow W^\pm W^\pm \rightarrow \ell^\pm \nu \ell'^\pm \nu'$ channel at 8 TeV [91–93]. This channel is advantageous at the LHC, because, contrary to all other VV states, the EW-induced production dominates the QCD-induced production. The first measurement of VBS production of a massive VV state with statistical significance exceeding 3.0 standard deviations (σ) is presented by the ATLAS Collaboration in the analysis of Ref. [92]. The EW WZ production was also studied at 8 TeV by the ATLAS Collaboration, which places

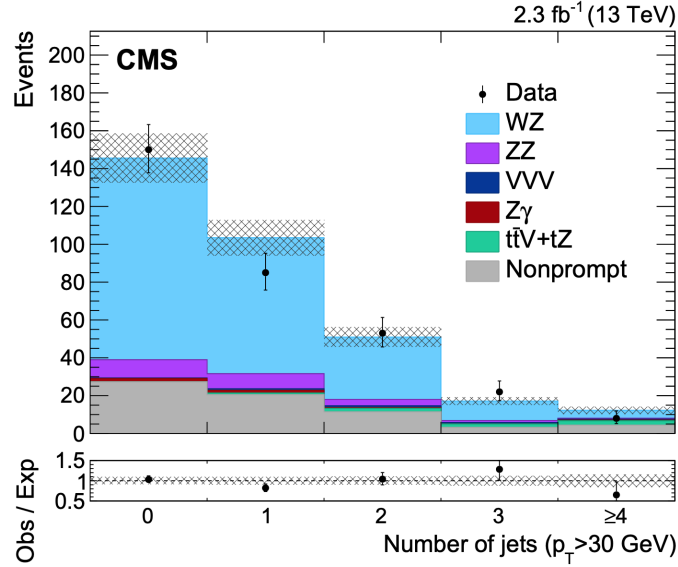


Figure 2.8: The number of jets in selected $\ell'\ell'\ell\nu$ events for the observed data and the expected WZ signal and background contributions using 2.3 fb^{-1} of data collected by the CMS experiment in 2015 [87].

upper limits on the production cross section in Ref. [84], and the CMS Collaboration, which reports a WZjj cross section in a region sensitive to EW WZ production in Ref. [91].

Studies of EW $W^\pm W^\pm$ have recently been performed by the CMS and ATLAS Collaborations at 13 TeV [94, 95]. The first measurement of EW VV production with a statistical significance greater than 5 standard deviations, considered the standard for discovery of new processes by the experimental particle physics community, is presented by the CMS Collaboration in Ref. [96]. A subsequent study performed by the ATLAS Collaboration [95] reports a comparable observation for the process. The first study of the EW VV production with a Z boson was performed in the $pp \rightarrow ZZ$ channel by the CMS Collaboration [97]. The observed statistical significance of the measurement is 2.7σ with 1.9σ expected in the SM.

The analysis of this thesis, submitted for publication to the journal *Physics Letters B* [55], is the first study of EW WZ production performed by the CMS. An analysis

using a comparable data set of pp collisions at 13 TeV has also been submitted to the journal *Physics Letters B* by the ATLAS collaboration [98]. The ATLAS analysis reports and observation of the EW WZ process with observed statistical significance of 5.3σ , with 3.4σ expected in the SM.

2.6.3 Searches for anomalous quartic gauge couplings

Anomalous triple and quartic gauge couplings were studied with measurements of WW and $WW\gamma$ production in electron-positron collisions at the LEP Collider at CERN by the ALEPH, DELPHI, L3, and OPAL Collaborations [99]. These processes are sensitive to the triple vector boson couplings WWZ and $WW\gamma$, and the quartic couplings $WW\gamma\gamma$ and $WWZ\gamma$. The production rate and angular distributions of the decay products of the vector bosons in these states were used to derive values of the triple and quartic couplings. No deviations were observed from the predictions of the SM were observed. Limits on the anomalous triple gauge coupling WWZ [100] were also presented by CDF and D0 collaboration in Ref. [80, 101], followed by studies performed by the ATLAS and CMS Collaborations in Refs. [84, 88].

Experimental sensitivity to the quartic coupling WWZZ was first achieved at the LHC. In pp collisions, quartic WZ interactions are accessible through VBS VV production or through triple vector boson production (VVV). Constraints on the deviations of quartic couplings from the SM prediction in WZ events with at least two jets were first presented by the ATLAS collaboration at 8 TeV [84]. This thesis reports the first study of aQGCs in the WZ channel at 13 TeV.

Fig. 2.9 illustrates the constraints on the operator parameters and scale of new physics for the dimension-8 quartic operators $\mathcal{L}_{T_i}$ from previous studies by the ATLAS and CMS Collaborations. The most sensitive measurements are those performed in the $W^\pm W^\pm$ channel reported in Ref. [94] and in the ZZ channel, reported in Ref. [97].

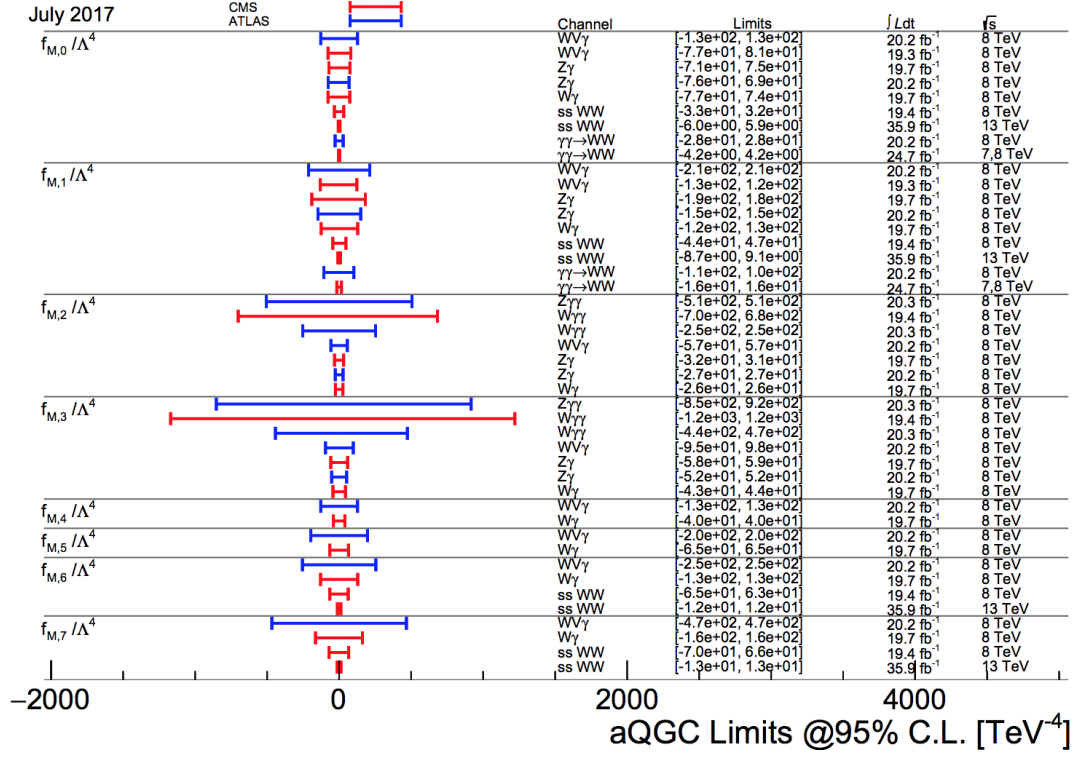


Figure 2.9: Constraints on the parameters f_{Mi}/Λ^4 for dimension-8 EFT operators from the CMS and ATLAS experiments.

2.6.4 Searches for charged Higgs boson production via vector boson fusion

The first studies of the production of charged Higgs bosons produced via VBF and decaying to WZ at the LHC were performed by the ATLAS Collaboration at 8 TeV [102]. The analysis exploited the decay channel $W^\pm \rightarrow q\bar{q}'$, $Z \rightarrow \ell\ell$ to place the first direct production constraints on the GM model, as well as model-independent limits on the production of a narrow-width charged Higgs produced via VBF. The CMS Collaboration performed the first study of the VBF $H^\pm$ exploiting the leptonic decay channel $WZ \rightarrow \ell\ell\ell'\nu$, using 15.2 fb^{-1} of data collected at 13 TeV [103]. The ATLAS Collaboration recently performed an analysis of the same channel using the data delivered by the LHC in 2016, corresponding to 36.1 fb^{-1} [104], extending the constraints of the

previous CMS analysis. A local excess of events over the SM prediction at a resonance mass of around 450 GeV was suggested by this analysis, with a global significance of 1.9σ when interpreted in the GM model. The analysis presented in this thesis extends the results of the previous CMS study, and is comparable to the recent study by the ATLAS Collaboration.

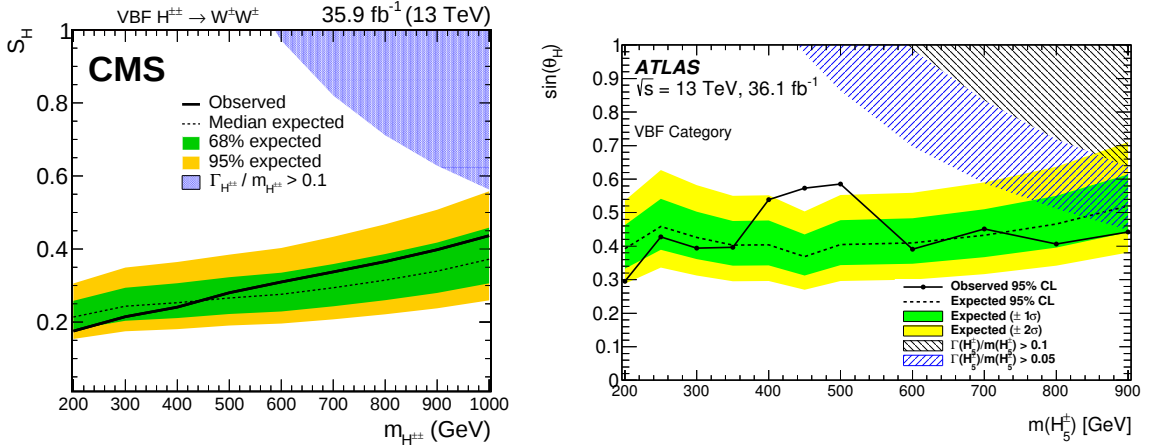


Figure 2.10: Limits on production of doubly charged (left, from Ref. [94]) and charged Higgs bosons (right, from Ref. [104]) in the H5plane defined by the GM model.

Members of the custodial fiveplet in the GM model have a common mass m_{H^5} . Therefore, limits on doubly charged Higgs bosons $H^{\pm\pm}$ are complementary to those obtained for $H^\pm$ production when interpreted under this model. The ATLAS and CMS Collaborations have studied $H^{\pm\pm}$ production with decays to same-sign W bosons at 8 TeV [91, 92], and the CMS Collaboration has presented constraints at 13 TeV [94]. The constraints on the VBF production of $H^\pm$ in the GM model from Ref. [104] and for $H^{\pm\pm}$ from Ref. [94] are shown in Fig. 2.10. The limits are presented in terms of s_H^2 and the mass of the $H^\pm$ or $H^{\pm\pm}$ boson, which defines the “H5plane” proposed by the LHC Higgs Cross Sections Working Group in Ref. [105]. The H5plane is defined with the assumptions that the triplet mass is greater than the fiveplet mass and that the parameters of the model allow perturbative calculations. Additional constraints on

the GM model arise from studies of b hadron decay and heavy Higgs boson searches, which are impacted by contributions from additional Higgs bosons contributing in loop diagrams, but these constraints are exceeded by the direct measurements discussed here. Global fits of the direct and indirect constraints on the GM model can be found in Ref. [53].

2.7 Motivation and content of this result

The work presented in this thesis has two related motivations: performing a measurement of a rare process predicted by the SM, and searching for signs of BSM physics in a channel and topology sensitive to modifications of the EW sector of the SM. Characterizing the self-interactions of the vector bosons is an important step to understanding the self-consistency of the SM, and a useful probe of possible deviations from its predictions. Because the production of events with multiple vector bosons in the final state requires collisions with a high center of mass energy, the LHC has the potential to explore these states in more depth than previously achieved.

We perform complementary interpretations of the observed data and its consistency with the prediction of the SM or signs of new physics. In particular, we measure

- The total production rate of $WZjj \rightarrow 3\ell\nu jj$ in fiducial regions with enhanced contributions from EW WZ production
- The consistency of the EW WZ production rate with the SM prediction and the statistical significance at which this rate can be established as non-zero.

We also place constraints on hypothetical new physics using the kinematic distributions of selected WZjj events with a VBS topology. These processes would be characterized by an excess of VBS WZ events with high m_{WZ} , due to the modifica-

tions of the quartic VVVV interaction in a manner sensitive to the scattering energy.

We place limits on the production and model parameters of

- Charged Higgs bosons, in a model-independent manner (assuming a small $H^\pm$ width) and in the Georgi-Machacek model.
- Anomalous quartic gauge couplings, using the language of dimension-eight effective field theory.

Chapter 3

The CMS Experiment at the CERN LHC

The Large Hadron Collider (LHC) at the European Organization for Nuclear Research (CERN) is the world's largest and most powerful particle collider [106]. It collides protons at a center of mass energy $\sqrt{s} = 13$ TeV, and lead ions at 2.76 TeV per nucleon. The protons or ions are brought to collision at the center of four large detectors, themselves a combination of many detector technologies. The particle detectors at the LHC were conceived, developed, and constructed over decades by collaborations of thousands of scientists to achieve a broad program of research. The detectors were designed and are operated by mutually exclusive collaborations of researchers from nearly all nations of the world. Collaborations of researchers maintaining and operating each detector at the LHC work independently, which enables cross validation of results while fostering productive competition. The detector sites along the LHC tunnel are:

- A Large Ion Collider Experiment (ALICE) detector [107], located at access point 2. Designed for the study of lead-ion collisions, especially for the charac-

terization of quark gluon plasmas formed in collisions.

- A Large Toroidal LHC ApparatuS (ATLAS) detector [108], located at access point 1. A general-purpose detector designed for sensitivity to new physics at the electroweak scale decaying to stable SM particles, in particular the discovery and characterization of the SM Higgs boson.
- The Compact Muon Solenoid (CMS) detector [109], located at access point 5. A general purpose detector with comparable measurement potential and research goals to the ATLAS detector.
- The Large Hadron Collider Beauty (LHCb) detector [110], located at access point 8. Designed to characterize the production and decay of b-quark hadrons. Particularly focused on charge–parity violations in these decays.

Results in this thesis are based on an analysis of proton-proton (pp) collisions at the LHC collected by the CMS detector in 2016. This chapter describes the design principles of the LHC and the CMS detector, as well as their operating characteristics in 2016.

3.1 The Large Hadron Collider

The primary design considerations of a particle collider are the energy of the particle collisions, the rate of collisions, and the objects collided. The collided objects determine the possible interactions that are accessible, whereas the energy of the collision drives the mass scale which is probed, subject to the constraint that the mass of any produced particle be less than the center of mass energy of the collision. As particle interactions are quantum mechanical and fundamentally stochastic in nature, the rate of collisions is also critical for achieving statistically significant measurements.

Several characteristics of pp collisions motivated the decision to use protons in the world’s highest-energy collisions. Accelerating and directing particles to collision relies on electric and magnetic forces, so the collided particles must be electrically charged. The only charged, stable, unbound, and fundamental particle is the electron. Electron–positron collisions have played a major role in the field of particle physics, however, because accelerated charged objects lose energy proportional to m^{-4} via synchrotron radiation, accelerating electrons to the TeV scale would require a prohibitively large tunnel to limit the bending force of the accelerator. The mass of the proton is nearly two thousand times greater than the electron mass, so synchrotron radiation is much less severe for accelerated protons. Unlike the electron, the proton is a composite object, so the conditions of its constituent quarks and gluons cannot be directly controlled when directing the protons to interaction, rather, they must be inferred from the properties of the colliding protons. These unknowns, governed by the complex structure of the proton, make pp collisions inherently messier than collisions of fundamental particles. Nonetheless, the compositeness of the proton is also advantageous due to the rich phenomenology of interactions proceeding via quarks and gluons. Additionally, quark–antiquark pairs are accessible in the proton at high energy scales, so interactions proceeding via particle–antiparticle annihilation are still accessible in high energy pp interactions. Without the need to produce antiparticles before bringing them to collision, as was necessary at lower-energy accelerators—notably the Large Electron Positron Collider (LEP) collider at CERN and the Tevatron collider at Fermilab in the United States—the LHC is able to deliver significantly higher interaction rate to complement its high collision energy [111].

3.1.1 Design characteristics of the LHC

The LHC was proposed to the CERN council, the management group of the laboratory, in 1994, and accepted with a preliminary budget in 1995. Construction was initially completed in 2008, with full operation beginning in 2009 after setbacks encountered in the initial commissioning. The collider is located on the outskirts of Geneva, Switzerland, which has been the site of the CERN laboratory since it was founded in 1954. The project was funded by the 20 member states of CERN, as well as monetary and research contributions from many other nations participating in the project, including the United States [112].

The LHC is situated in a 26.7-km-circumference tunnel, 45–170 m below ground. The tunnel pre-dates the LHC, having been originally built for the Large Electron Positron (LEP) collider. It consists of 8 straight sections 528 m in length, and 8 arced sections. The cross section of the majority of the tunnel is 3.7 m in diameter, with larger excavated areas at the four experimental caverns and other access points. The primary motivation for an underground tunnel is to circumvent the high cost of land acquisition, but underground operation is also advantageous for reducing the cosmic radiation reaching the experimental cavern and for shielding the radiation produced by the LHC [106]. Protons are directed through pipes inside the tunnel, which are held at high vacuum. The positions and accelerations of the protons are controlled by magnetic and electric fields maintained by instrumentation surrounding the vacuum pipes. The arced sections are equipped with dipole magnets, which direct charged objects along a circular path. The size of the LEP tunnel prohibited the installation of two independent beam systems, which led to the adoption of a unique “two-in-one” superconducting magnet design [113], shown in Fig. 3.1.

Neglecting synchrotron radiation, the primary limitations on the energy of a circular collider come from the magnetic field of the bending magnets and the radius of

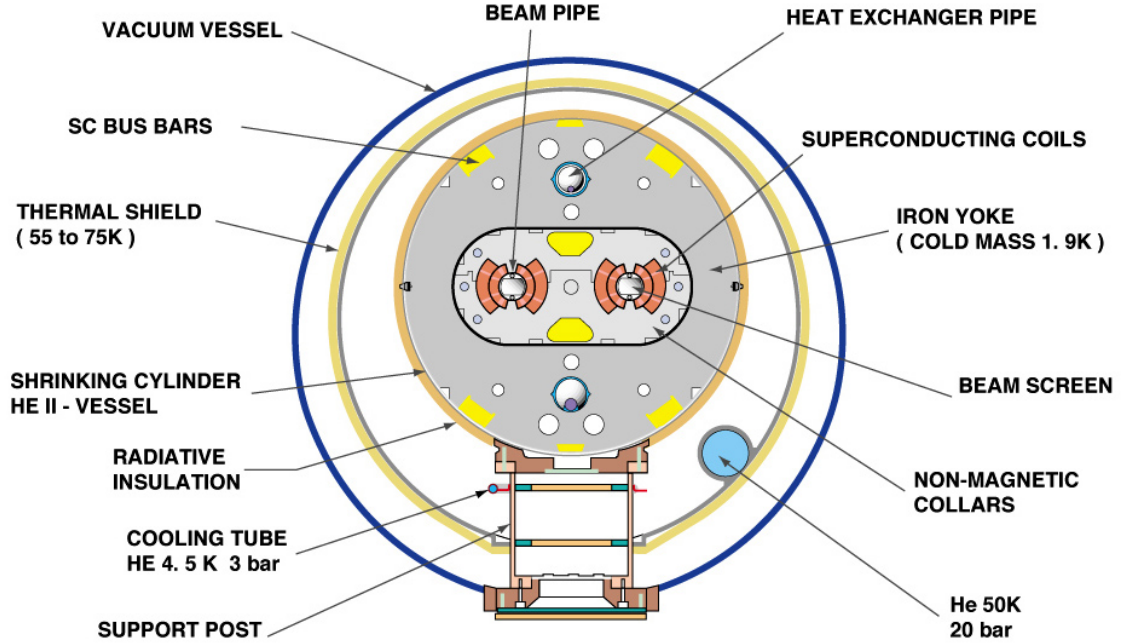


Figure 3.1: Cross section of an LHC dipole magnet [114].

curvature of the tunnel. A particle of charge q with velocity v will travel through a magnetic field of strength B along a circular trajectory. These quantities are related to the energy of the particle and its radius of curvature R by

$$E = qBRv \approx qBRc, \quad (3.1)$$

where $v \approx c$, the speed of light, for $p \gg m$. Increasing the radius requires a larger tunnel, which is limited by the cost of construction. The LHC tunnel size was set by the existing LEP tunnel, so achieving a sufficient magnetic field was a major focus of the LHC development.

The LHC dipole magnets, with a maximum field strength of 8.33 T, were constructed at an affordable cost through major technological advances. Such a high field necessitates the use of superconducting magnets. The magnets are constructed of niobium-titanium, and are operated at 1.9 K, lower than any previous accelerator. Because about 80% of the arced sectors of the LHC are equipped with dipole

magnets, the bending radius of the dipoles is 2.804 km, yielding the design energy of 7 TeV per beam from Equation 3.1. In addition to dipole magnets, superconducting quadrupole, octopole, and sextapole magnets are used to constrain the beam into its compact form.

The protons used in the LHC are acquired by exposing hydrogen gas to a strong electric field to ionizes the hydrogen atoms, leaving free protons. The protons are not injected directly into the LHC ring, but are instead accelerated through a chain of linear and circular accelerators of increasing size. This acceleration chain, in the context of the full accelerator complex at CERN, is shown in Fig. 3.2. The protons are accelerated from rest to an energy of 50 MeV in the linear accelerator LINAC2 before being fed into the Proton Synchrotron Booster, where they reach an energy of 1.4 GeV. The Proton Synchrotron Booster feeds the Proton Synchrotron (PS), the world's highest-energy collider in 1959, where the protons reach 25 GeV. The PS feeds the Super Proton Synchrotron (SPS), where the protons reach 450 GeV before injection into the LHC.

The protons are accelerated from injection energy to the collision energy by radiofrequency (RF) cavities housed in the straight section of the LHC tunnel at access point 4. The protons are injected into the LHC in bunches, which must be matched to the 400 MHz oscillating frequency of the electric field of the RF cavities. Protons are injected from the SPS into the LHC such that every 10th RF oscillation is populated, leading to the LHC collision rate of 40 MHz. The field oscillations leads to stronger or weaker kick applied to in-time or out-of-time protons, which works to further coalesce the beam into bunches. The 25 ns bunch separation corresponds to a total of 3564 possible bunches that can be stored in the LHC. However, to avoid damaging the LHC cryogenic system when removing the proton beam from the LHC, it is necessary to direct the beam away from its nominal path to deposit its energy in

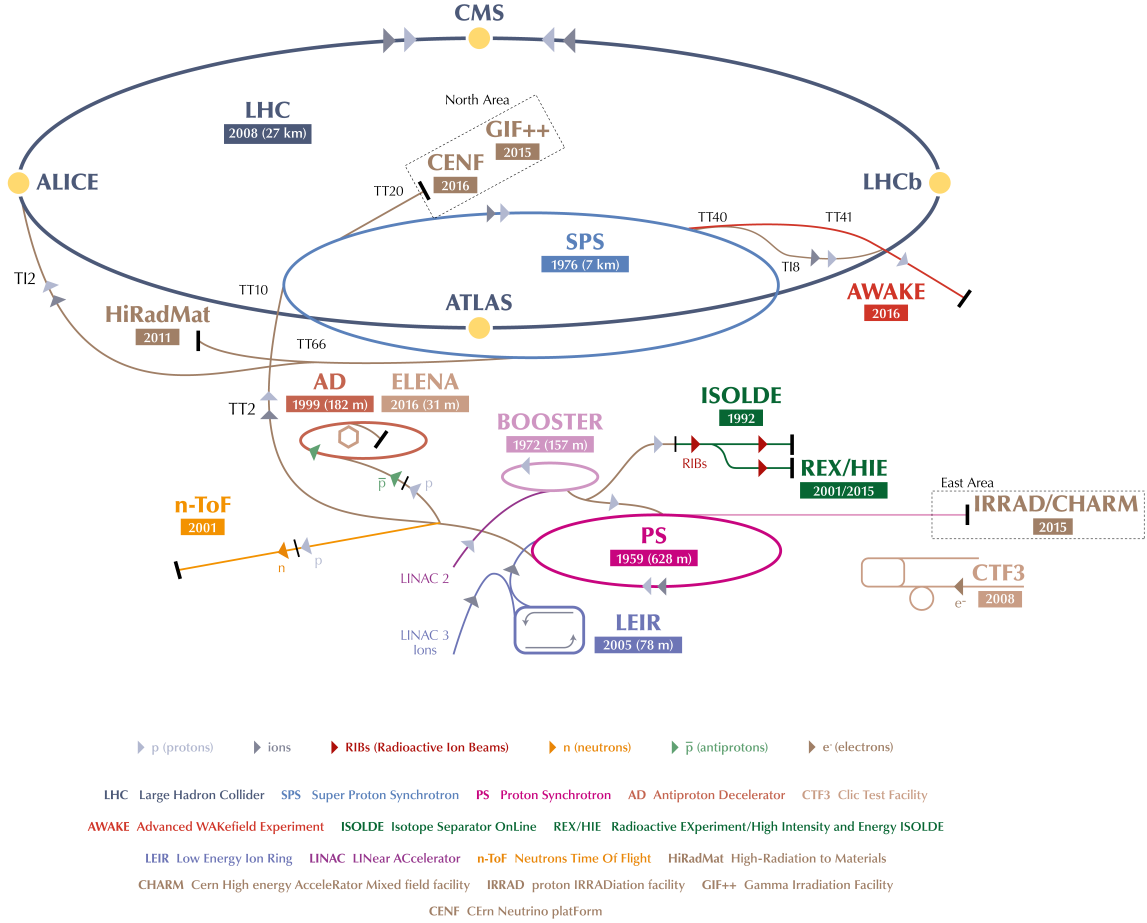


Figure 3.2: The accelerator complex at CERN, including the LHC and its feeder accelerator chain [115].

a controlled way. This process, known as a beam dump, is facilitated by kicker magnets in access point 6 which deflect the beam. A gap of protons is necessary to allow the magnet time to reach the necessary field without scanning the beam through the cryogenic system, so only a maximum of 2808 bunches are ever populated by protons.

3.1.2 Operation of the LHC in 2016

In 2016, the LHC delivered pp collision at $\sqrt{s} = 13$ TeV from April 22 to October 27. During this time, the proton beams were circulating in the LHC for about 75% of the

time, that majority of which was devoted to delivering collisions for physics analysis. The LHC maintained a very high level of performance, even exceeding its design parameters by some metrics. However, operation at $\sqrt{s} = 13$ TeV is below the design energy of $\sqrt{s} = 14$ TeV. Operating at a lower energy was necessary due to the time-consuming process of “training” magnets to hold a stable field without quenching, or sudden loss of superconductivity, which can be damaging to the magnet but also helps to align the magnetic coils, increasing the maximum field which the magnet can sustain. This procedure will again be undertaken during the LHC shutdown from 2019–2020, and collisions at 14 TeV are envisioned beginning in 2021.

The instantaneous luminosity $\mathcal{L}$ measures the rate of pp collisions delivered by the LHC and governs the rate of particle production in collisions. The maximum instantaneous luminosity achieved by the LHC in 2016 was $1.5 \times 10^{34} \text{ cm}^{-2} \text{ s}^{-1}$, which exceeds the design luminosity of $\approx 1.1 \times 10^{34} \text{ cm}^{-2} \text{ s}^{-1}$. The integrated luminosity $\mathcal{L}_{int}$ is obtained by integrating the instantaneous luminosity over time. The total number of collisions of interest is directly proportional to the integrated luminosity, so maximizing both the active time devoted to collisions and the instantaneous luminosity are vital to achieving the LHC physics program. Integrated luminosity is measured in cm^{-2} or b^{-1} , where the barn (b) is defined as $1 \text{ b} = 10^{-24} \text{ cm}^2$. The LHC data set is generally reported in fb^{-1} .

The instantaneous luminosity is purely a function of the LHC beam parameters. Defined in terms of the number of protons per bunch N_b (assuming symmetric beams), the number of proton bunches n_b , the beam revolution frequency f_{rev} and the effective area of beam overlap A_{eff} , it can be expressed as

$$\mathcal{L} = \frac{n_b N_b^2 f_{rev}}{A_{\text{eff}}} . \quad (3.2)$$

Assuming the beams have Gaussian profiles, and introducing the variables β^* , the

amplitude of the betatron function at the interaction point, the normalized emittance ϵ_N , a measure of the beam spread in position and momentum space, the beam crossing angle θ_c , and σ_z and σ^* , the longitudinal and transverse RMS widths of the bunches in the lab frame, the expression becomes

$$\mathcal{L} = f_{rev} \frac{n_b N_b^2 \gamma}{4\pi \beta^* \epsilon_N} \left(1 + \left(\frac{\theta_c \sigma_z}{2\sigma^*} \right)^2 \right)^{-1/2}, \quad (3.3)$$

Peak values of these parameters from 2016 running are shown in Table 3.1.

Table 3.1: Peak performance parameters contributing to LHC luminosity in 2016 operation [116].

Parameter	2016 value	
Protons per bunch (N_b)	$\approx 1.1 \times 10^{11}$	
Number of bunches (n_b)	2220	
Normalized emittance (ϵ_N)	≈ 2.0	mm·mrad
β^*	40	cm
Crossing angle (θ_c)	280	μrad
Bunch length	≈ 1.05	ns
Peak luminosity	≈ 1.5	$\times 10^{34} \text{ cm}^{-2} \text{ s}^{-1}$
Peak mean pile-up	≈ 45	

Because of the high density of protons when the beams are brought to collision, there are generally multiple pp interactions per beam crossing, referred to as pileup. Pileup interactions increase the multiplicity of particles in the detector, many of which are not associated with the interaction of interest. The number of pileup interactions per collision is driven by the instantaneous luminosity and the inelastic pp cross section, $\sigma_{pp} \approx 70 \text{ mb}$. An average of 23 pp interactions per event were observed in the 2016 data set. Techniques to mitigate the impact of pileup when selecting events of interest are discussed in Chapter 5.

3.2 The Compact Muon Solenoid experiment

The CMS detector is a general-purpose detector designed to study particle production and interactions at the TeV scale. A major design principle of the CMS detector was the ability to probe the nature of electroweak symmetry breaking, which has been achieved through the discovery and characterisation of a scalar boson consistent with the SM Higgs boson [35, 36]. The design of the CMS detector ensured sensitivity to a Higgs boson with mass up to 1 TeV in its SM decays, as well as sensitivity to a broad class of BSM theories. In terms of detector performance, these design principles, as defined in Ref. [109], are:

- Accurate muon identification and high momentum resolution, including dimuon mass resolution of $\approx 1\%$ at $m_{\mu\mu} = 100$ GeV and correct charge assignment up to $p_T^\mu \approx 1$ TeV.
- Identification of objects with a short but significant decay time, including tau leptons and b hadrons, which requires fine position resolution close to the interaction point to distinguish their displaced tracks.
- Good electromagnetic energy and momentum resolution, achieving diphoton and dielectron mass resolution of $\approx 1\%$ at $m_{ee} = 100$ GeV, and sufficient granularity to distinguish prompt diphotons from π^0 decays to photons.
- Sufficient calorimeter resolution and hermeticity for precise dijet mass and missing transverse momentum reconstruction.

The CMS detector has a cylindrical geometry. As shown in Fig. 3.3, it is built from the combination of several detector technologies which work in harmony to achieve the outlined design goals. The following sections briefly describe each of the detector

subsystems which comprise the CMS detector. Chapter 5 details the ways in which these subsystems are used together to reconstruct the physics objects used for the results presented in this thesis.

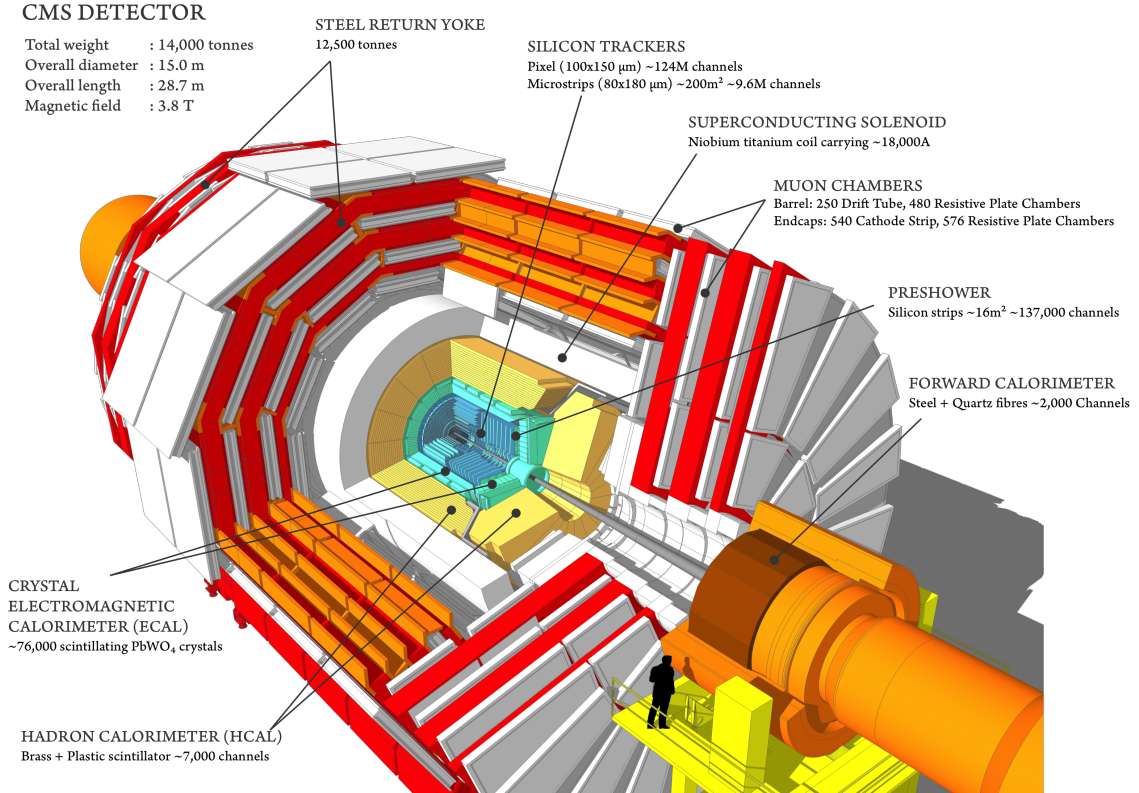


Figure 3.3: A cutaway view of the CMS detector, showing each detector subsystem [117].

The coordinate system adopted by CMS, and used in this thesis, is a right-handed coordinate system with the point $(x, y, z) = (0, 0, 0)$ centered in the detector, at the nominal collision point. The y -axis is perpendicular to the earth, where vertically upward defines the $+y$ direction. The x -axis is in the plane of the LHC, with the $+x$ direction pointing towards the center of the ring. The z -axis points through the center of the detector along the direction of travel of the colliding protons, where $+z$ is defined by the right-handedness of the coordinate system. In polar coordinates, the azimuthal angle ϕ is measured from the x -axis in the x - y plane, and the polar

angle θ is measured from the z -axis. The radial coordinate r is defined as $r = \sqrt{x^2 + y^2 + z^2}$. Because the production of particles is preferentially in the forward direction (along the z -axis), it is convenient to introduce the pseudorapidity $\eta = -\ln(\tan \theta/2)$ for the polar coordinate. The momentum in the transverse direction, defined as $p_T = \sqrt{p_x^2 + p_y^2}$, is a particularly important quantity, because the initial transverse momentum of the collision is zero. In this thesis, the four-momentum of an object will commonly be expressed as $p = p(m, \vec{p}) = p(m, p_T, \eta, \phi)$.

3.2.1 The CMS solenoidal magnet

The central feature of the CMS apparatus is a superconducting solenoid of 6 m internal diameter and 12.5 m in length, constructed from 4 layers of reinforced niobium titanium. It provides a nearly constant 3.8 T magnetic field inside the solenoidal volume. The magnetic flux of the solenoid is returned through a 12000-tonne steel yoke, comprising 5 wheels and 2 endcaps, which is fully saturated to approximately 2 T. The tracker and calorimeters are situated inside the solenoid, while the muon detectors are embedded in the steel return yoke, as shown in Fig. 3.3.

During operation, the CMS solenoid stores approximately 2.4 GJ of energy, the largest magnet in the world by this metric. This immense size leads to a powerful bending radius for charged objects within the detector, which is a critical metric for particle reconstruction and identification. In particular, objects of charge q moving in a magnetic field $\vec{B}$ at velocity $\vec{v}$ experience a Lorentz force,

$$\vec{F} = q\vec{B} \times \vec{v}. \quad (3.4)$$

Because the solenoidal field is constant and aligned along the z -direction, $\vec{B} = B\hat{z}$, the force is purely in the transverse direction. Consequently, charged particles within the CMS solenoid travel along a nearly helical path of radius $R = p_T/|q|B$, where

small deviations from this path arise from non-uniformity of the field and interactions in the detector material. It follows that the p_T of a particle of known charge can be deduced by measuring its radius of curvature, therefore, the bending power of the magnet BR is a critical parameter determining the detector capability for accurate particle p_T measurement. The nearly 12 Tm bending power achieved by the CMS is central to achieving the high momentum resolution outlined in the design goals of the experiment. It is also fundamental to the particle flow reconstruction technique utilized by CMS, outlined in Chapter 5.

3.2.2 The CMS silicon tracking system

The inner part of the CMS detector is instrumented with semiconductor tracking detectors. The tracking system is positioned inside the solenoidal magnet, surrounding the interaction point, with a cylindrical shape of length 5.8 m and diameter 2.5 m. Sensors arranged in discs cover the forward part of the volume ($\eta < 2.5$). It consists of silicon pixels in the innermost part, and silicon strips further from the interaction point. The tracker geometry, including the layout of the silicon sensors, is shown in Fig. 3.4.

The tracking system is designed to measure the helical paths of charged particles emanating from the collision point, from which the transverse momenta can be inferred. Charged particles passing through the reverse-biased doped silicon diodes produce ionization currents, which are measured to infer the position of the interaction with very high position resolution. With 200 m² of active silicon area, the CMS tracker is the largest silicon tracker ever built.

At the operating conditions of the LHC in 2016, an average of 23 proton interactions took place per bunch crossing, producing on the order of 10³ particles every 25 ns. In this environment, the tracking system is crucial to determining which par-

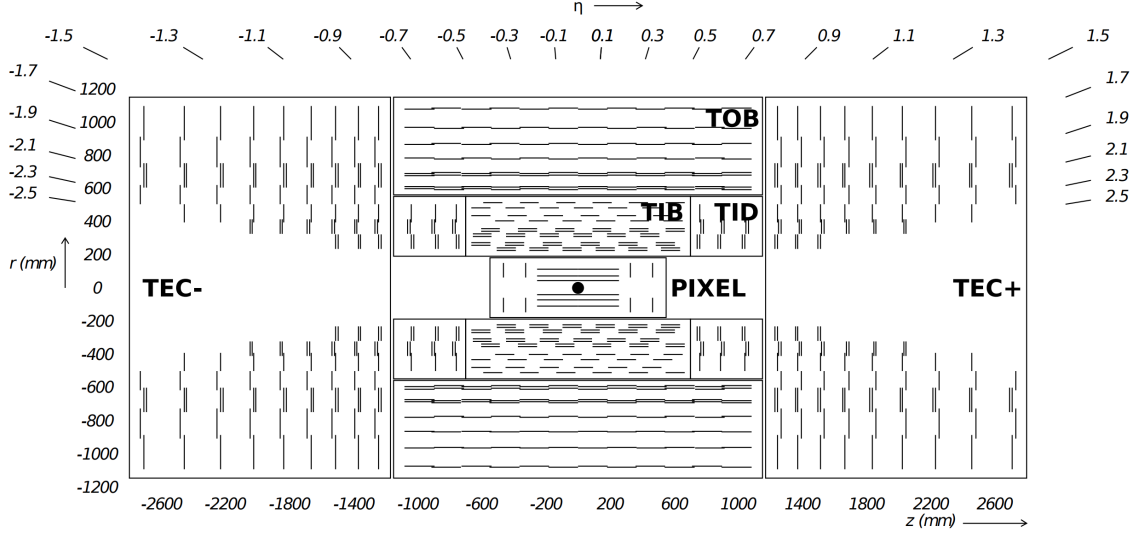


Figure 3.4: Cross section of the CMS tracker, viewed in the r - η plane. At the center is the pixel detector, with the tracker inner (TIB) and outer barrel (TOB) instrumenting the barrel part of the detector at increasing distances from the beam line. The tracker inner disk (TID) and tracker endcap (TEC) are perpendicular to the beam pipe. Single lines represents a detector module, and double lines represent double-sided modules. Reproduced from Ref. [118].

ticles are associated with the interaction of interest. In addition, decays displaced from the production vertex is an important signature of heavy-flavor hadrons and τ leptons, so knowledge of the track origins in three dimensions is crucial. The paths of charged particles can be deduced from their discrete interactions in the tracking system, from which the particle origin can be inferred. Because multiple interactions in a single sensor complicate track reconstruction, the pixel and strip dimensions are designed to keep the particle flux per collision low. The pixel detector barrel consists of three layers, positioned at 4.4, 7.3, and 10.2 cm from the beam line. At the inner layer, the hit rate density is $\approx 1 \text{ MHz/mm}^2$, which motivates the use of $100 \times 150 \mu\text{m}$ silicon sensors, for an occupancy on the order of 10^{-4} per sensor per bunch crossing. The flux is significantly reduced further from the beamline, so the outer parts of the tracker uses silicon strip detectors, 10-25 cm in length and $80\text{-}180 \mu\text{m}$ in length, to

reduce the channel count and cost. In total, the silicon tracking system has over 75 million read-out channels. The position resolution of individual sensors is $10 - 40 \mu\text{m}$, allowing for vertex association which is accurate to $\approx 10 \mu\text{m}$, and momentum resolution for muons with $p_{\text{T}} = 100 \text{ GeV}$ and $\eta < 1.5$ is around 2% [119].

3.2.3 The CMS electromagnetic calorimeter

Immediately outside the tracking system, but inside the solenoid, is the electromagnetic calorimeter (ECAL). It is constructed from transparent lead tungstate (PbWO_4) crystals, which absorb energy of electromagnetically-interacting objects (e.g., electrons and photons) before emitting light of wavelength $\lambda = 420 - 430 \text{ nm}$, with intensity proportional to the absorbed energy. The energy deposits of the interacting particles are obtained from photodetectors that collect the emitted light. Lead tungstate was chosen for properties including its high density of 8.3 g/cm^3 , which contributes to its small radiation length—the mean distance an electron travels in the material before losing $1/e$ of its initial energy—of 0.89 cm , and its small Molière radius (the radius at which 90% of the electromagnetic shower is contained on average) of 2.2 cm . These properties facilitate the compact design of the CMS ECAL. Because photons and electrons interact significantly with the large physical structure of the CMS solenoid, positioning the ECAL outside of the solenoid would degrade the energy performance and compromise the association of tracks with ECAL energy deposits. Consequently, a compact calorimetry system is required to reduce the volume—and therefore the cost—of the solenoid. In addition, the scintillation decay time of the crystals, the delay between the absorption of energy and light emission, is such that 80% of the light is released by an excited crystal within 25 ns , enabling the necessary fast read out. These highly favorable characteristics of PbWO_4 are partially offset by its relatively low light output of 4.5 photons per MeV of deposited energy at the operating

temperature of 18° C.

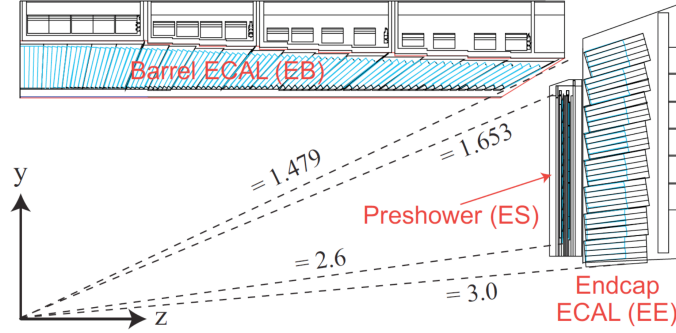


Figure 3.5: An r - η cross section of the CMS ECAL, showing the detector components, crystals, and their η coverage. Reproduced from Ref. [120].

The ECAL system consists of 61,200 PbWO_4 crystals in the barrel, $|\eta| < 1.444$, and 7,324 crystals in each of two encap discs covering $1.566 < |\eta| < 3.0$. The small gap in the ECAL is due to cabling and support structure for the tracking system. The cross section of the crystals in the plane facing the interaction point is $2.2 \times 2.2 \text{ cm}^2$, equal to the Molière radius of the material. The barrel crystals are 23 cm in length, whereas the endcap crystals are 22 cm, corresponding to 25.8 and 24.7 radiation lengths respectively. The arrangement of the crystals is shown in Fig. 3.5. Because of the relatively low light yield of PbWO_4 , fast photodetectors with powerful amplification capabilities, that are capable of operating in the high magnetic field, are needed. Avalanche photodiodes instrument the barrel, while vacuum phototriodes are used in the barrel, due to their superior radiation hardness.

The resolution of the energy measurement in the ECAL improves with energy up to 0.5–1 TeV, at which point the percentage of the shower energy not contained in the ECAL becomes significant. Below this threshold, higher energies correspond with greater light production, and a more accurate energy measurement. The uncertainty

in the energy measurement below ≈ 500 GeV can be parameterized as

$$\left(\frac{\sigma_E}{E}\right)^2 = \left(\frac{S}{\sqrt{E/\text{GeV}}}\right)^2 + \left(\frac{N}{E/\text{GeV}}\right)^2 + C^2. \quad (3.5)$$

Where S is a stochastic term, driven by the number of scintillation photons produced, N is a noise term arising from the electronics and background radiation produced outside the primary pp interaction, and C is a constant term due to the non-uniformity of the crystal material and fabrication. The values of $S = 2.8\%$, $N = 0.12$ and $C = 0.30\%$ were measured in the initial CMS commissioning [109], and are illustrative of the performance during 2016 data taking.

3.2.4 The CMS hadronic calorimeter

Outside the ECAL, but inside the solenoid, is the hadronic calorimeter (HCAL), covering the radial space from 1.77 m to 2.95 m. The HCAL measures energy of charged and neutral hadrons for identification and measurement of hadronic jets. The hermetic coverage of the HCAL up to $|\eta| < 5.0$ permits an extensive characterization of the transverse energy of hadronic particles in an event, allowing neutrinos (or weakly interacting new physics) to be detected as an unbalance in transverse momentum. Unlike the ECAL, the HCAL is a sampling calorimeter, that is, it is constructed of layers of dense material to induce hadronic showers interspersed amongst active layers of scintillating material. The barrel ($|\eta| < 1.3$) and endcap ($1.3 < |\eta| < 3.0$) region of the HCAL are constructed from flat steel and brass absorber plates, 50-80 mm thick, layered between scintillating tiles. The scintillators are segmented to cover an area in (η, ϕ) of $(0.087, 0.087)$ for $|\eta| < 1.6$ and $\approx (0.17, 0.17)$ for $|\eta| > 1.6$. Wavelength shifting fibers adjust the wavelength to a more favorable frequency before light is collected and measured by hybrid photodiodes. At $\eta = 0$, the HCAL is $5.82\lambda_I$ thick, whereas it is $\approx 10\lambda_I$ thick at $|\eta| > 1.3$. Here λ_I is the nuclear interaction length,

the mean free path of a hadronic particle before undergoing an inelastic hadronic interaction. Hadronic particles encounter about $1.1\lambda_I$ of material before reaching the HCAL.

In order to measure very forward jets—which are critical for this result—a forward HCAL system (HF) is installed close to the beam pipe on the extreme end of the detector, ± 11.15 m from the interaction point. The HF detects particles in the region $3.0|\eta| < 5.0$. Because the energy deposit per collision in the HF is 5–10 times greater than the combined energy deposits in the rest of the detector, the HF must be extremely radiation hard. The active material of the HF is quartz fiber Cherenkov detectors, which collect radiation from particles traveling with velocity above the Cherenkov limit. The fibers are embedded in dense steel to induce electromagnetic and hadronic showers. Half of the fibers extend the full length of the detector, whereas half begin at a depth of 22 cm, which gives sensitivity to the shower development, to distinguish between hadronic and electromagnetic showers. The detectors are housed in a hermetic steel, concrete, and polyethylene shielding of 0.85 m in thickness for additional radiation hardness. The resolution of the HF is at least a factor of two lower than the barrel and endcap calorimeters. Despite this, it plays an important role in the results presented in this thesis, which rely on forward jet reconstruction.

The layout of the HCAL, showing each subsystem discussed, as well as the small outer part of the HCAL outside the magnet, is shown in Fig. 3.6. As for ECAL, the energy resolution can be described by Equation 3.5. The results obtained from pion test beams before the start of CMS data taking are $S = 115\%$ and $C = 5.5\%$. The noise term N is subdominant.

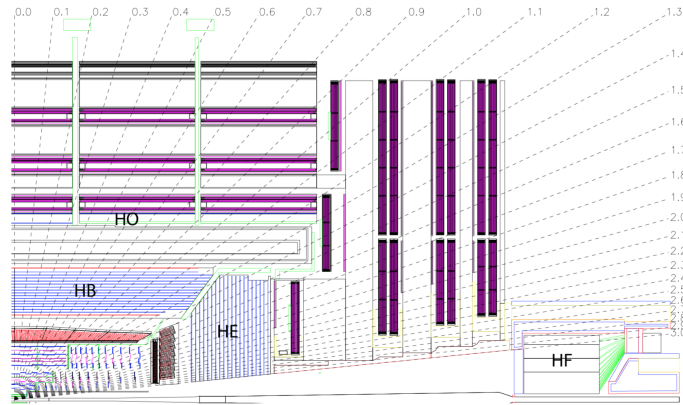


Figure 3.6: An r - η cross section of a quadrant of the CMS detector. The components of the HCAL system are shown. [35].

3.2.5 The CMS muon system

As suggested by the CMS name itself, the muon detector subsystem is essential to the achieving the experimental goals of the CMS experiment. High-energy muons are produced in many heavy-object SM and BSM decays, and long-established detector technologies are well suited to muon detection in LHC collisions. Muons with $v \approx c$ are nearly stable on timescale of the traversal of the CMS detector volume. As the heaviest stable lepton, they are less susceptible to electromagnetic radiative energy loss than electrons, and therefore interact minimally with the calorimeter system. Detector technologies suitable for charged-particle detection, placed outside all other detector systems, are therefore ideal for muon detection.

The muon system is composed of three technologies, described in more detail in the following sections. Each operates on the principle of gas ionization. The detectors consist of a volume of gas over which an electric field is applied. Charged particles passing through the volume ionize the gas, and the ionized gas molecules and free electrons drift in the electric field towards read-out electronics at the edges of the detector volume. The track of the interacting object is inferred from the position and timing of the measured electrons and ions. The operating characteristics of the gas

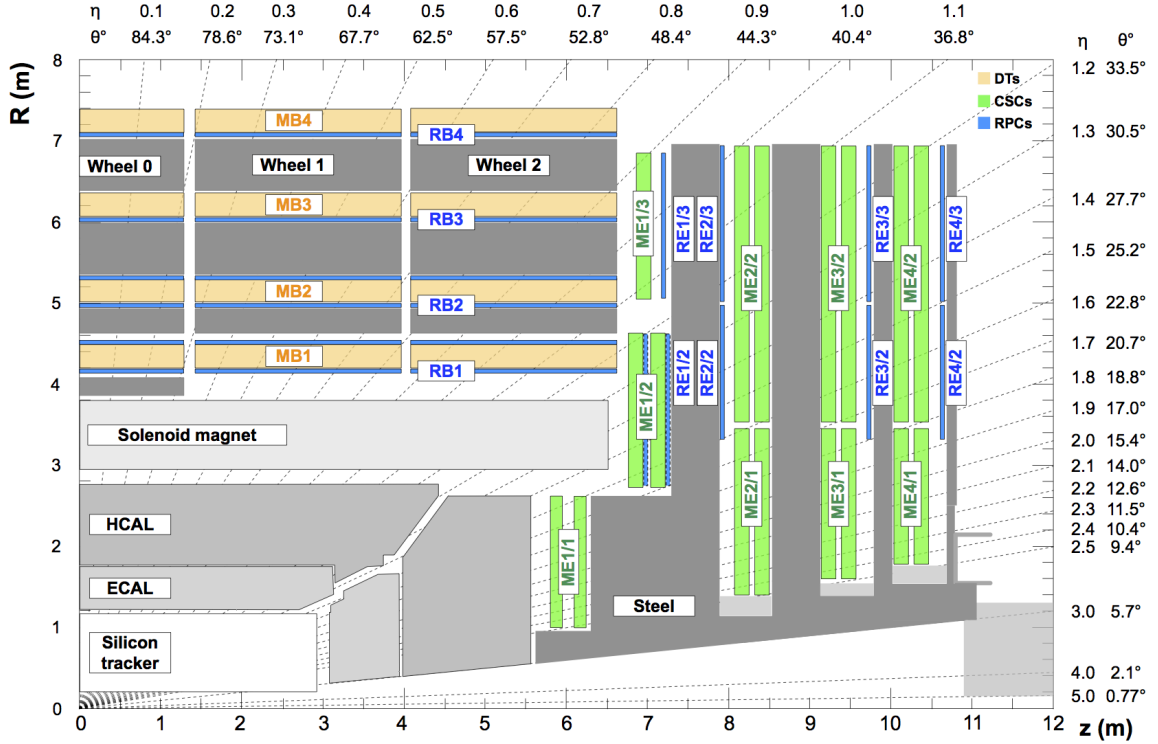


Figure 3.7: An r - η cross section of a quadrant of the CMS detector, with the interaction point at the bottom left corner. The locations of the detectors comprising the CMS muon system within the steel steel return yoke are shown [35].

ionization detectors can be modified for fast readout, high granularity, and In the CMS detector, these considerations, along with the cost of construction for detectors covering an area of $\approx 25,000 \text{ m}^2$ and the effect of the radiation environment of the detector on performance, motivate the specific design choices.

Gas ionization detectors are sensitive to all charged particles passing through the detector. As previously discussed, and illustrated in Fig. 3.7, the muon system is the outer layer of the CMS detector, embedded in the steel return yoke of the magnet system. This structure, along with the inner detector systems, serves as shielding for the muon system and ensures that the well over 99% of objects reaching the muon system are muons.

3.2.5.1 Drift tube system

The barrel part of the muon detector, $|\eta| < 1.2$, is composed of drift tube (DT) chambers. In this region, the magnetic field outside the return yoke is generally small (below 0.2 T), and the muon occupancy is low, justifying the use of DT detectors. The DT system is built from individual DTs, rectangular aluminum tubes with a length of 2.4 m and a cross section of $13 \times 42 \text{ mm}^2$, shown in Fig. 3.8. At the center of each tube is a $50\text{-}\mu\text{m}$ -diameter gold-plated stainless-steel anode wire. Aluminum strips are attached to the interior of the aluminum structure—insulated by mylar tape—to shape the electric field to the pattern illustrated in Fig. 3.8, and an 85% Ar 15% CO_2 gas mixture is contained in the volume. The drift time for muons crossing a cell at a maximum distance from the anode wire of 21 mm is about $380 \mu\text{s}$, which is sufficient to keep the occupancy low. Muons passing through the gas volume ionize the gas and create an avalanche of electrons around the very high field at the anode wire. A time-to-digital converter reads out the electrical pulse of the electrons. The arrival time of the signal, corrected by the drift velocity and the time from the collisions to the trigger decision, are used to obtain the muon interaction position.

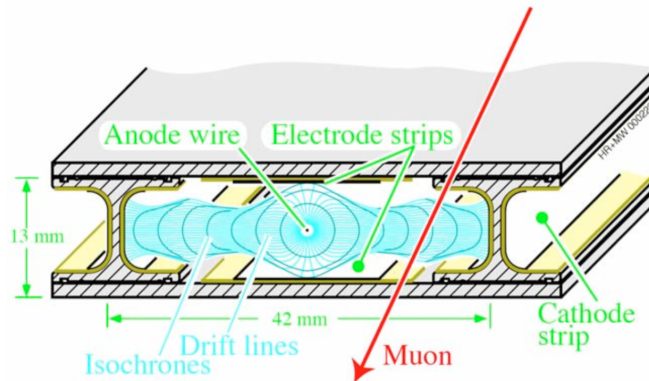


Figure 3.8: Diagram of a drift tube cell, from Ref. [109]. The geometry of the cell and the electric field lines resulting from the anode wire and electrode strips are illustrated.

The DT cells are assembled into superlayers (SL), composed of 4 layers of drift cells staggered by half a cell. The SLs are grouped into chambers of 2 or 3 SLs which are arranged into stations within the magnet return yoke. The wires in the outer SLs within a chamber are parallel to beam line, providing a measurement in the bending plane. An additional inner SL is aligned perpendicular to the beam direction, for all chambers but those in the outermost station, to provide a measurement in the z plane. The chambers are arranged in concentric cylinders around the beam line, with 60 chambers each in the three inner cylinders and 70 in the outer cylinder. In total there are $\approx 172,000$ cells, read out as independent channels. The offset arrangement of the cells eliminates dead spots in the efficiency, and allows a measurement of the muon crossing time to be obtained using meantimer circuits. The muon position resolution in the DTs is around $100\,\mu\text{m}$ and the timing resolution is within a few ns.

3.2.5.2 Cathode strip chamber system

Cathode strip chambers (CSC) are used in the forward regions of the CMS detector ($0.9 < |\eta| < 2.4$), where the prompt and background muon rates are high, and the magnetic field is large and non-uniform. The CSC system is constructed of trapezoidal multi-wire proportional chambers, which subtend 20 or 30 degrees in ϕ . The largest chambers are 3.4 m long and up to 1.5 m wide. The chambers are divided into 6 sections—7 mm wide for the innermost layer of chambers and 9.5 mm for all others—each filled with an Ar-CO₂-CF₄ gas mixture. Aluminum cathode strips, which span the radial direction at a constant phi, are milled to 6 of the panels forming the sections. The strips vary from 8.4 to 16 mm, for a constant width in ϕ , and are separated by about 0.5 mm. In addition, each layer has a plane of anode wires running lengthwise, $50\,\mu\text{m}$ in diameter, separated by 2.5 mm, and slightly tilted to correct for the deflection of the drifting electrons in magnetic field. Muons passing

through the chambers create an avalanche of electrons in each layer, which is read out on the anode wires. In addition, a mirror pulse is created in the cathode strips. Together, the pulses and shapes give an excellent position measurement in both the $r - \phi$ and η planes, as shown in Fig. 3.9. Each cathode strip is read out, whereas anode strips are read out in groups of 16. The position resolution of a CSC chamber is around $100\ \mu\text{m}$, and the timing resolution is $\approx 7\ \text{ns}$

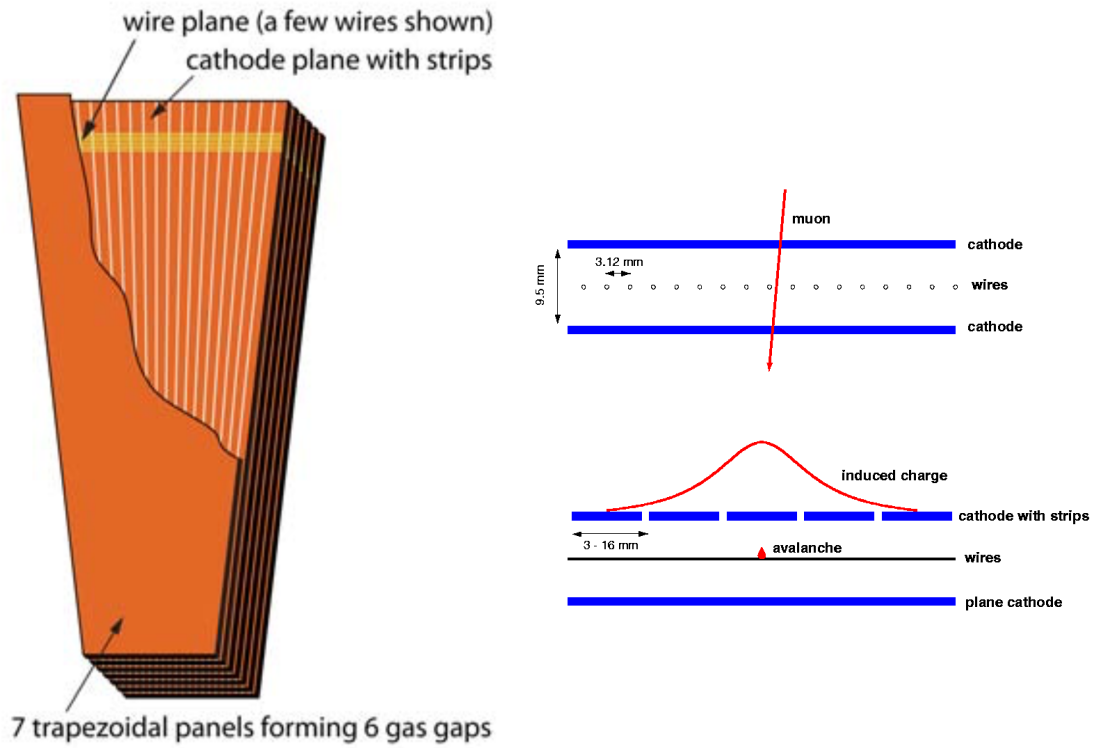


Figure 3.9: Left: diagram of a CSC chamber. Right: principle of operation for a CSC layer, showing the avalanche around the anode wires and the pulse shape on the cathode strips, which are combined for accurate position measurement. Both figures are reproduced from Ref. [109]

The CSC chambers are arranged perpendicular to the beam line, conjoined for full coverage in ϕ . Four stations of CSC chambers are interspersed amongst the magnetic flux return plates in each endcap of the detector, see Fig. 3.3. The innermost station (for each endcap) is composed of 3 rings of CSC in the radial direction, whereas other

stations have two rings. In total there are 540 CSC chambers. All but the outer ring of chambers in the innermost layer partially overlap the neighboring chambers, avoiding regions of low efficiency. Members of the University of Wisconsin – Madison CMS group have played a crucial role in the design and operation of the CSC system since the inception of the CMS experiment. The author of this work contributed to the operation of the CSC system, serving the detector-on-call for the detector subsystem for three weeks per year from 2016 – 2018. During these periods the author represented the CSC system in daily meetings, coordinated maintenance and development projects, and responded to critical issues preventing data collection or degrading data quality.

Aging studies of the CSC chambers at the original CERN Gamma Irradiation Facility (GIF), and at the new GIF++, have shown that the CSC system performance is resistant to the damage incurred by the radioactive environment of high luminosity collisions. The author contributed to the on-going studies at the GIF++ facility by monitoring the physical degradation of CSC chambers exposed to radiation, including changes to the leakage current and gas gain with increasing exposure. The continued performance of the CSC chambers after exposure to sustained radiation, comparable to many times the luminosity already delivered to CMS, has provided evidence that the CSC system will remain an important part of CMS throughout the full LHC program.

3.2.5.3 Resistive plate chamber system

A crucial role of the muon system is to provide muon position and momentum measurements in a sufficient timescale to determine if an event justifies storage offline analysis. This process, known as event triggering, is described in additional detail in Section 3.2.6. While the DT and CSC system provide fast readout, the background

rate and importance of associating an event with the correct bunch crossing (within 25 ns) make additional redundancy desirable.

A system of resistive plate chambers (RPC) in both the barrel and endcap ($|\eta| < 1.2$) accomplishes this goal. The RPCs are double-gap chambers constructed from a thin layer of readout strips between two electrodes held at high voltage, forming a volume filled predominately with $\text{C}_2\text{H}_2\text{F}_4$ gas. The system is operated in avalanche mode, where the sum of the signals in the two gaps created by an interacting muon is read out on strips at the center plate. The plates are separated by 2 mm, which ensures that the charge avalanche is observed well below the 25 ns window needed to assign the muon to the correct bunch crossing. There are 6 layers of RPCs in the detector barrel, interspersed among the DTs, and 3 layers in the endcap alongside the CSCs. The readout strips in the barrel RPCs subtend an angle of $5/16$ degree. The position resolution of the RPCs, about 1 cm, is below the DT and CSC systems, but the system provides an independent position and timing measurement.

3.2.6 Event triggering and data acquisition

Despite significant zero-suppression and data compression, the read out from the millions of channels in the CMS detector for a typical pp collision of interest requires around 1 MB of storage. Transmitting this quantity of information at the collision rate of 40 MHz delivered by the LHC to the CMS detector exceeds the threshold of practical technological capabilities. Furthermore, the event rate of rare standard model processes and theorized BSM interactions is many orders of magnitude below this rate. Therefore, it is neither feasible nor desirable to store every collision.

The CMS trigger system uses a two-tiered approach to select events of interest—typically characterized by having one or more objects at high transverse momentum—and to reject more common pp interactions, such as elastic electromagnetic scattering.

The Level 1 (L1) trigger is a custom-designed hardware system, primarily built from programmable electronics, that is designed to reduce the rate of collisions to below 100 kHz in a latency of about $4\,\mu\text{s}$. The L1 trigger decision, known as an “L1 accept” when an event is stored, uses coarsely segmented information from the muon and calorimeter systems. Information in the inner tracker is not used at L1, due to the read-out latency of the system.

The L1 trigger system initially handles the calorimeter and muon system information independently. For the calorimeter information, the ECAL crystals are grouped into towers of 25 crystals, to match the HCAL granularity in η and ϕ space and to reduce granularity. Energy deposits in the ECAL and HCAL are used to build primitive objects, e.g., the energy and shape of ECAL deposits contribute to photon and electron candidates while HCAL information suggests clusters of hadronic particles produced in the event. The particle candidates are sorted by p_{T} , and the total energy deposit in the event is computed, before forwarding the event information to the global trigger for an event-wide decision.

The muon trigger system builds track candidates by matching hit patterns in the CSC, the DT, and the RPC subsystems. Track finders specific to the barrel region ($|\eta| < 0.9$), the endcap region ($0.9 < |\eta| < 1.2$), and the CSC-DT overlap region ($1.2 < |\eta| < 2.4$), combine the track segments from individual detectors to form muon candidates. The candidates built by each track finder are combined and disambiguated by the global muon trigger, which sorts the candidate objects by p_{T} before forwarding them to the global trigger.

The global trigger uses the sorted candidate and event information from the muon and calorimeter triggers to test physic objects and object combinations, based on kinematic and quality criteria, before accepting or rejecting the event. The set of conditions by which events are evaluated are organized into “trigger paths,” which

act to sort the dataset and the procedure by which it is evaluated. For paths with very high rate, e.g., signal lepton triggers with lower p_T thresholds, a trigger path may be “prescaled.” This process reduces the event rate by rejecting an unbiased subset of the events otherwise satisfying the acceptance criteria. Prescaled trigger paths are important for validating the efficiency and performance of other trigger paths.

Data collected by the detector system is continually read and stored into 40-MHz pipelined buffers while waiting for a trigger decision to determine if the data should be overwritten by the continual stream of new data, or pushed to the next stage of the data acquisition system for storage. The $4\mu s$ window for the trigger decision is set by the storage capacity of the most limited buffer system, associated with the silicon tracker. If the L1 trigger determines an event should be read for further analysis, the data are forwarded over a high-bandwidth network system to a processing farm for further analysis. The read-out bandwidth limits the rate of accepted events at L1 to about 100 kHz. In some cases, it is necessary to explicitly enforce this condition, accomplished through the trigger throttling system. The trigger throttling system prevents event read out in cases where a subsystem data buffer is near capacity, or when event triggering in quick succession would lead to excessive rate. For example, the trigger throttling system prevents event read out for three successive bunch crossing after a triggered event. This reduction in event readout, referred to as deadtime, is generally unbiased and kept below 1%. However, for this analysis, a systematic loss of events due to mistimed trigger signals lead to an inefficiency in collection of events of interest. This effect is discussed in detail in Chapter 6.

Events accepted by the L1 trigger undergo an additional filtering by the high level trigger (HLT) before they are stored to tape for offline analysis. The HLT is a software-based system running on a commercial processing farm that has access to the full event information from all detector systems. In principle, the HLT can

implement the same reconstruction algorithms described in Chapter 5, but the time constraints for the decision and the limited size of the compute farm make this impractical. Instead, the HLT implements similar algorithms with optimizations and simplifications designed to reach an event acceptance or rejection decision in less than $200\,\mu\text{s}$. As with the L1 trigger, the event evaluation is divided into trigger paths. Operations are performed per trigger path such that uninteresting events are rejected as soon as possible, with computationally expensive operations such as track building only performed at the final stages. The high cost of large-scale data storage, especially in a way that allows quick access for processing at sites throughout the world, is an increasing concern for the LHC experimental program. In addition to ensuring manageable read-out rates, efficient filtering by the HLT is an important aspect of addressing this storage issue.

3.2.7 Luminosity measurement

The luminosity delivered by the LHC to the CMS interaction point is a critical input parameter to any cross section measurement and many LHC searches, as it is needed to establish the expected background contribution from processes described by MC simulation and to convert an observed event rate to a measured cross section. Luminosity is defined purely in terms of the parameters of the LHC, as given by Equation 3.3. However, because the conditions of the LHC beams are not constant throughout a data collection period (e.g., the number of protons in the colliding bunches decreases continually during an LHC fill), the instantaneous luminosity must be continually monitored. Establishing the luminosity in terms of the parameters in Equation 3.3 while the LHC is delivering collisions is unfavorable because measurements of the beam parameters, such as the beam widths, cannot be made to high precision without interfering with the beam conditions. Instead, the luminosity can be

monitored by exploiting the relationship between luminosity and detector occupancy, which is linear for the current for certain detector systems and under the current LHC operating conditions [121]. Therefore, the relationship between the observed occupancy rate and the luminosity can be expressed as

$$R = \sigma_{vis}\mathcal{L}, \quad (3.6)$$

where the proportionality constant σ_{vis} is known as the visible cross section seen by the detector system. By establishing $\mathcal{L}$ and σ_{vis} in conditions where the parameters of Equation 3.3 are known, further monitoring of the luminosity can be achieved purely via a measurements of the R .

The reference luminosity measurement is accomplished in dedicated LHC run periods using a technique established by Simon Van der Meer at the CERN intersecting storage rings in 1968 [122]. During these periods, the x and y positions of the LHC beams are scanned from their nominal head-on alignment used during collision by controlled distances. Because the number of colliding protons depends on the beam overlap, the occupancy rate in the detector can be used to establish the Gaussian profiles of the beams. Combined with the rate at nominal collision alignment and the beam intensity, this procedure, referred to as a Van der Meer scan, establishes the relationship of Equation 3.6. The CMS experiment monitors the luminosity by measuring the occupancy in the pixel detector, the DT system, and the HF. In addition, two detector systems dedicated to the luminosity measurement are used, the fast beam conditions monitor and the pixel luminosity telescope. A separate R_i and $\sigma_{vis,i}$ is obtained for each monitoring system, so the consistency of the independent measurements can be used for validation. The consistency of different approaches to monitoring the luminosity, along with uncertainties in establishing the beam parameters in the Van der Meer scan, contribute to the systematic uncertainty of the

luminosity measurement. For the 2016 dataset, the luminosity is established with a systematic uncertainty of 2.5% [123]. In 2016, the LHC delivered 41.0 fb^{-1} of pp collisions to the CMS interaction point.

3.2.8 Data quality

The CMS detector collects data from collisions delivered by the LHC to the CMS interaction point with all components of the CMS detector active, configured, and synchronized. In some cases, issues with a subsystem status or configuration prevent data collection. Reasons for data loss include, e.g., trigger throttling (discussed in the previous section), or problems with high or low voltage power supplies that require manual intervention to recover detector operation. In 2016, the CMS experiment collected 92% of the data delivered by the LHC.

Collected data are evaluated by experts from each subsystem to assess the detector performance, because problems that did not prevent data collection can still compromise its utility for physics analysis. Collected data with severe localized inefficiencies that prevent object reconstruction in a kinematic region and bias results in a complex way are excluded from the data set used for analysis. This analysis uses only data certified by the CMS experiment for physics analysis, which corresponds to 95% of the collected data.

I represented the CSC subsystem as the primary data certification expert during the collection of data used in this thesis. Because of the large redundancy of the CSC chamber layout, the CSC system approved over 99% of the collected data for physics analysis. A water leak at the end of 2016 in the CSC cooling system necessitating a large-scale shutdown of the inner CSC chambers was the major source of data loss, however, this primarily effected pp collisions at 5 TeV, which are not used in this thesis. The source of the problem was identified during the early-2017 LHC shutdown,

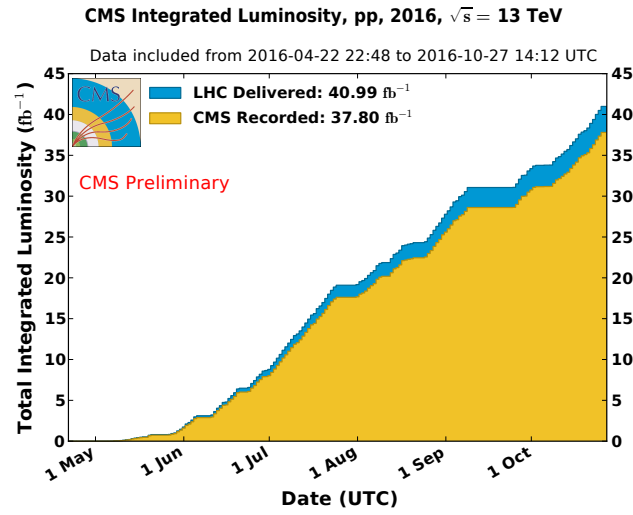


Figure 3.10: The accumulated luminosity of the data delivered to the CMS interaction point, collected by the CMS experiment, and approved for physics analysis by detector experts, are shown. Reproduced from Ref. [124]

and repaired by reinforcing the connections of the copper pipe cooling system.

Chapter 4

Theoretical predictions and event simulation

Quantifying the agreement of experimental observations with the SM or with possible BSM theories is a core goal of collider physics. While the SM is a nearly complete and profoundly successful theory, connecting theories that concern the interactions and excitations of quantum fields to the electrical signals induced in the CMS detector by pp collisions is highly non-trivial. In an ideal case, one would use the theory of particles and fields outlined in Chapters 1 and 2 as input to derive the expected distribution of measurable signals resulting from pp collisions. In this paradigm, BSM modifications to the SM are realized as modifications to the structure of the underlying QFT. They lead to new interactions, which modify the rate or type of particle production, or kinematics variables of the produced particles, leading to deviations from the SM expectation in measured quantities. In practice, a factorized approach—leveraging approximations and tuning to experimental data—allows the simulation of primary particle production in pp collisions, the formation of bound states and particle decay, and the interactions of particles in the CMS detector to be

simulated in a way that largely achieves this goal. The CMS Collaboration benefits from the work of collaborations of theoretical physicists and many previous studies to produced the detailed simulations used to model signal and background processes for this analysis.

Nonetheless, it is important to highlight that experimental physics is not purely at the mercy of the simulations from which predictions derive. Many predictions are phenomenological in nature, i.e., they have been tuned to the observations measured in LHC collisions and previous accelerators. Furthermore, in situ measurements allow the data observation to modify aspects of the predictions. Lastly, broad knowledge of the nature of particle production in collisions can sometimes be leveraged to completely remove dependence on simulation. For example, with no resonant source of diphoton production, the $m_{\gamma\gamma}$ distribution would be a falling spectrum. The first measurements of the scalar Higgs boson with $m_H = 125 \text{ GeV}$ made in the diphoton channel where achieved with the SM $pp \rightarrow \gamma\gamma$ expectation parameterized as a falling exponential distribution, not with an ab initio simulation of the distribution [125, 126].

The results in this thesis exploit the data to constrain important backgrounds. However, the EW WZ and QCD WZ processes are only distinguishable by their kinematic distributions, which cannot be independently observed. Determining the production mechanism of WZjj production and identifying the component that is sensitive to the WWZZ coupling relies heavily on achieving well-understood simulations of the signal and background processes contributing to the events analyzed. This chapter discusses the techniques used to simulate pp collisions at the LHC. The techniques used to obtain results for WZjj production in the SM and beyond, and the predictions used in this analysis, are presented and discussed.

4.1 Anatomy of an LHC collision

As discussed in Chapter 2, perturbation theory is the most effective technique for calculations of particle scattering in the SM QFT. While the underlying principles are well-established, calculations of QCD interactions, in particular, are computationally challenging. The QCD theory is asymptotically free, meaning $\alpha_s(\mu)$ grows with low values of the energy scale μ , so low-energy phenomena are non-perturbative. Because asymptotic freedom also leads to bound states at the low energy scales of most physical phenomena, the energy regime of stable particles—and therefore the incoming state directed to collision and the outgoing particles that interact with the detector—have non-perturbative properties.

The uncomfortable dichotomy between the regime of most physical phenomena and calculability is averted by the concept of factorization, which allows a separation of long-distance interactions from short-distance interactions, governed by energy scales small or large with respect to $\Lambda_{\text{QCD}} \approx 200 \text{ MeV}$, respectively. Factorization states that the interaction of hadrons can be reduced to structure functions describing the distribution of quarks and gluons in hadrons, known as parton distribution functions (PDF), convoluted with the parton-parton interaction cross section. At high momentum transfer, the parton-parton interaction can be described perturbatively. The evolution of the interaction from outgoing quarks and gluons to the formation of bound hadronic states, described by non-perturbative parton-to-hadron fragmentation functions or models, occurs at lower energy and larger time and length scales. Along with the decays of bound states, they are separable from the perturbative interactions and can be built from phenomenological models tuned to independent measurements. The interactions of stable particles with the detector material are separated from their production in time and space, allowing them to be treated inde-

pendently with sophisticated models of particle interactions in matter implemented in the GEANT4 package [127, 128]. The concept is illustrated for pp collisions in Fig. 4.1. Factorization is formally established under certain conditions [129], and it has been highly successful in describing experimental observations over generations of hadron-hadron, lepton-lepton, and hadron-lepton colliders.

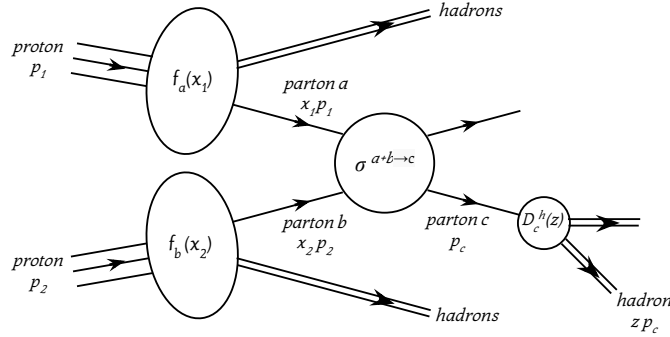


Figure 4.1: Illustration of the principle of QCD factorization. The cross section of the process $pp \rightarrow X$ is reduced to the parton distribution functions $f_{a,b}(x)$, the partonic cross sections $\sigma^{a+b \rightarrow c}$, and the parton-to-hadron fragmentation functions $D_c^h(z)$. Reproduced from Ref. [130].

Factorization allows the challenge of describing LHC collisions to be tackled in pieces, leveraging different techniques for each piece while benefiting from possibly independent experimental data. This chapter summarizes how these factorized components are modeled for the simulations that are used to guide in this analysis. In some cases, it is useful to consider only some contributions. For example, the effect of the experimental reconstruction can be parameterized by reconstruction efficiencies, or, in cases where the effects are not dramatic and are well-established, the hadronization of quarks and gluons can be reduced to an effective smearing. This features allow calculations that neglect aspects of the full simulation to still be used for meaningful comparisons with experimental measurements.

4.2 Parton distribution functions

A fundamental consequence of factorization is that the low-energy confinement of quarks and gluons in the proton can be described independently from the parton-parton interaction in a hard collision, e.g., momentum transfer $Q^2 \gg \Lambda_{\text{QCD}}^2$. Mathematically, the separation of the proton structure and the high- Q^2 parton interaction, shown pictorially in Fig. 4.1, can be expressed as

$$\sigma^{\text{pp} \rightarrow X} = \sum_{a,b \in \{q,g\}} \int dx_1 dx_2 f_a(x_1, \mu_F^2) f_b(x_2, \mu_F^2) \hat{\sigma}^{ab \rightarrow X}(x_1 x_2 s, \mu_F^2) \quad (4.1)$$

where x_1 and x_2 are the fractions of the proton momentum carried by partons $a, b \in \{q, g\}$ and the PDFs $f_{a,b}(x_i)$ give the probability of extracting the given parton with momentum fraction x_i from the proton. The total momentum must be divided amongst the constituents, i.e.,

$$\sum_i \int_0^1 dx x f_i(x, \mu_F^2) = 1 \quad (4.2)$$

Here μ_F is the factorization scale, an energy scale which represents the transition from the non-perturbative regime of the PDF and the perturbative high Q^2 interaction. Though it is often convenient to take $\mu_F^2 = Q^2$, μ_F is a free parameter that is an artifact of the truncated series in perturbation theory. As the order of the perturbation expansion considered increases, the μ_F dependence of the result is reduced.

At least to very good approximation, the PDFs can be considered universal functions. Therefore, they can be derived using independent measurements, such as deep inelastic scattering of lepton and proton beams. As shown in Fig. 4.2, a lepton incident on a proton target interacts with the quarks of the proton via a virtual γ or Z . By controlling the incident energy of the proton and lepton and measuring the outgoing particles, the abundance and momentum fraction of the proton constituents can be established.

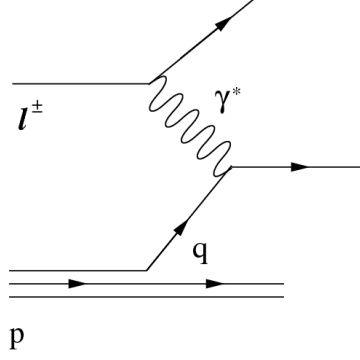


Figure 4.2: Illustration of deep inelastic scattering of a lepton from a proton, used as a probe of the internal parton structure of the proton. Reproduced from Ref. [131].

In practice, it is not feasible to parameterize all partons (and flavors) for all x_i through experimental measurements. Fortunately, while perturbation theory cannot predict the formation of bound states, it does allow a means to parameterize the evolution of the bound partons across energy scales. The scale dependence of the PDF is captured by the DGLAP equations, first established by Dokshitzer [132], Gribov [133], Altarelli and Parisi [134] in the 1970s:

$$\mu_F^2 \frac{df_a(x, \mu_F^2)}{d\mu_F^2} = \sum_{b \in \{q, g\}} \int_x^1 \frac{dz}{z} \frac{\alpha_s}{2\pi} \hat{P}_{ba}(z) f_b(x/z, \mu_F^2). \quad (4.3)$$

The functions $\hat{P}_{ba}(z)$ are the Altarelli–Parisi splitting functions, which describe the splitting of parton a into b , carrying a momentum fraction z of the initial momentum, that can be derived order by order in perturbation theory. The splitting of parton a is accompanied by an additional parton, e.g., $g \rightarrow q\bar{q}$, that is absorbed in to the parton sea of the proton. Given the DGLAP equations, if the proton PDF can be established for any complete phase space of parton flavors, momentum fractions, and Q^2 , it can be evolved to all other scales. In practice, the contributions to the PDF are measured at many different Q^2 with experiments sensitive to different quark flavors. The experimental data is fitted, considering the DGLAP evolution, to obtain a full

parametrization of the PDF at all energy scales.

Collaborations of theorists use independent techniques to fit data from many collider and fixed-target experiments to produce large data sets of predictions that can be used in simulations. The PDF is parameterized and distributed as data grids that can be sampled with a centralized interface known as LHAPDF [135]. Results in this thesis primarily make use of the NNPDF3.0 [136] set of PDFs, which use a neural net to generate pseudodata that is parameterized with global fits. Uncertainties in this procedure are assessed by evaluating the quality of the PDF fit to the data, comparing the predictions of independent collaborations in a formulaic way [137], and by determining the dependence of predictions on μ_F . The PDF uncertainty and its impact on the results of this thesis are further discussed in Chapter 6.

4.3 Perturbative calculations and matrix element generators

With the parton content of the proton captured by the PDFs, the next challenge in modeling an LHC collision is the perturbative calculation of the parton-parton interaction. As discussed in Chapter 2, perturbation theory relies on an expansion in the coupling constants of the theory,

$$\sigma = \sigma_0 \alpha_s^0 + \sigma_1 \alpha_s^1 + \sigma_2 \alpha_s^2 + \dots \quad (4.4)$$

For a process with only EW couplings at the lowest order, such as VV production, σ_0 is called the leading order (LO) cross section, and σ_1 the next-to-leading order (NLO). While α_s is sufficiently small at the LHC collision energy to justify the perturbative expansion ($\alpha_s(m_Z) = 0.1189 \pm 0.0010$ [49]), the energy and phase-space dependence of higher-order terms is non-trivial, leading to QCD corrections much larger

than the naive expectation in many cases [89, 138, 139]. The majority of production processes at the LHC require a calculation at least to NLO in QCD for percent-level accuracy. Theoretical tools have advanced to allow automated computation of all SM processes at NLO [140–142], and many processes have recently been computed at NNLO in QCD [62] or N³LO [143, 144]. Calculations at LO in the EW theory are often sufficient for percent-level accuracy. However, some processes are susceptible to anomalously large EW corrections, including VBS VV production [145], considered in this thesis.

The cross section of Equation 4.4 can be expressed in terms of the square of the scattering matrix element $\mathcal{M}$, which expresses the transition probability from the initial to the final state. The matrix element, or scattering amplitude, is calculable in perturbation theory, of which techniques based on Feynman diagrams are the most familiar. Because the total cross section is a scan over all possible outgoing energy and momenta combinations of the outgoing particles, the cross section is an integration over a many-dimensional phase space of configurations Φ_n ,

$$\sigma_i = \int d\Phi_n |\mathcal{M}_i|^2 . \quad (4.5)$$

The integration can be performed with numerical techniques. The high dimensionality of the phase space, as well as the complex peak structure arising from resonances, is well-suited for Monte Carlo (MC) integration techniques, described in the following section.

In the Feynman diagram approach to matrix-element calculations, amplitudes are represented by Feynman diagrams where QCD (EW) coupling vertices in the diagrams are proportional to $\sqrt{\alpha_s}$ ($\sqrt{\alpha}$). A calculation at a given perturbative order involves all diagrams connecting desired initial and final states for which the product of vertices is less than the order of the perturbative expansion. Feynman diagrams

can be generated in an automated procedure with techniques from graph theory, before being used to derive and calculate amplitudes from Feynman rules in the SM. Automation of this process was established in the early 1990s at LO [146]. Extending this to procedure to NLO has been accomplished more recently by the program `MADGRAPH5_aMC@NLO` [141], the successor to the previous work, which is used extensively for MC simulations in this thesis.

Diagrams without loop contributions are called “tree-level” diagrams. The lowest tree-level contribution to a given state, shown for $pp \rightarrow WZ$ in Fig. 4.3 (far left), is referred to as the Born-level contribution to the process. The NLO correction to the Born process includes contributions from tree-level diagrams of $pp \rightarrow WZ + j$ for $j \in \{g, q\}$, known as real emission diagrams, and loop contributions, shown in Fig. 4.3 (center) and (left) respectively. Divergences from each of these contributions cancel through renormalization. The delicate procedure of “cancelling infinities” presents a major challenge to automating NLO calculations with numerical techniques. It requires a cut to be made very close to the infinite pole, and careful management of the divergent contributions around the pole with high-precision numerical accuracy. All possible one-loop amplitudes in the SM have been calculated with analytic techniques. Automated codes access these loop contributions and residual divergences from packages such as `OpenLoops` [147], and combine them with the Born and real-emission calculations, while verifying that the infinities of the loop and real-emission calculations do indeed cancel.

For the first NLO calculations that became available, the interplay between contributions was adjusted per process in dedicated calculations. This procedure is now automated in several general-purpose MC programs used in this analysis, including `MADGRAPH5_aMC@NLO` [141], `SHERPA` [140], `HERWIG` [148], and `Recola` [142]. However, programs which automate all SM calculations are not optimized for specific

process, which can lead to prohibitively long computation time for high-multiplicity states or processes with a complex phase, such as EW WZ. For this reason, the best-available calculations for the QCD WZ and EW WZ process are remain those from dedicated implementations, such as Refs. [149, 150].

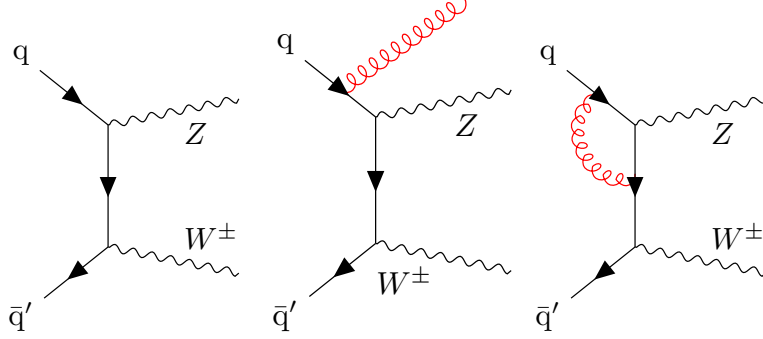


Figure 4.3: Contributions to the $pp \rightarrow WZ$ process at NLO. The leftmost diagram is the Born contribution, which is part of the LO calculation. The real-emission diagram is shown in the center, and the loop contribution is far right. The cross section is proportional to the square of the sum of amplitudes shown, which at $\mathcal{O}(\alpha_s)$ includes the square of the Born and real-emission diagrams and the interference of the loop and Born diagrams.

Diboson processes inclusive in the number of jets have recently become available at next-to-next-to-leading order (NNLO) [62]. These calculations have an additional layer of complication due to the more complex cancelling of divergences between double loop, double real-emission, and combined real-emission and loop contributions. They are not yet automated, partially because not all two-loop amplitudes in the SM are known. Combining these predictions with hadronization functions, discussed in the following sections, is also not yet accomplished for diboson processes. In this thesis, we use NNLO cross section calculations to correct the predicted yields of NLO simulations for diboson processes.

4.4 Monte Carlo integration and event unweighting

As referenced in the previous section, calculations of physical quantities involve integrations of the amplitude over the relevant phase space for the scattering process. The phase space integration in Equation 4.5 is an integration over the quantum numbers and four-momenta of all particles considered. The integration is often complex and rarely solvable analytically. The MC integration algorithm was first formally developed in the late 1940s as part of the Manhattan project, with the name derived from the eponymous casino for due to its leveraging of random processes [151]. It is well suited for the phase space integration, because the the uncertainty in the result scales independently of the dimensionality of the integration [152].

The MC procedure estimates the integral of a function $f(x_i, \dots, x_n \equiv \vec{x})$, defined over a volume V^n , by tossing random random points inside a known volume over which the function is defined, of dimension V^{n+1} . The volume V^{n+1} from which the random numbers are drawn must fully contain $f(\vec{x})$, however, no further knowledge of $f(\vec{x})$ over V^{n+1} is required. The function is evaluated for the coordinates $x_i \dots x_n$ of the random point. If the random point is inside the volume of $f(\vec{x})$, it is accepted, if it falls in a region of V^{n+1} outside of $f(\vec{x})$ it is rejected. As the number of points sampled N grows, the ratio of the number of accepted points to N times the volume of V^{n+1} approaches the integral of $f(\vec{x})$. This is illustrated for one dimension in Fig. 4.4. The error of the integration decreases proportional to $\sqrt{N}$, which is very poor for low dimension integration ($\lesssim 4$) but much better than traditional techniques for high-dimensional problems.

In addition to predictions for aggregated quantities such as total cross sections and differential rates, event-wise simulations are critical to modeling expected results

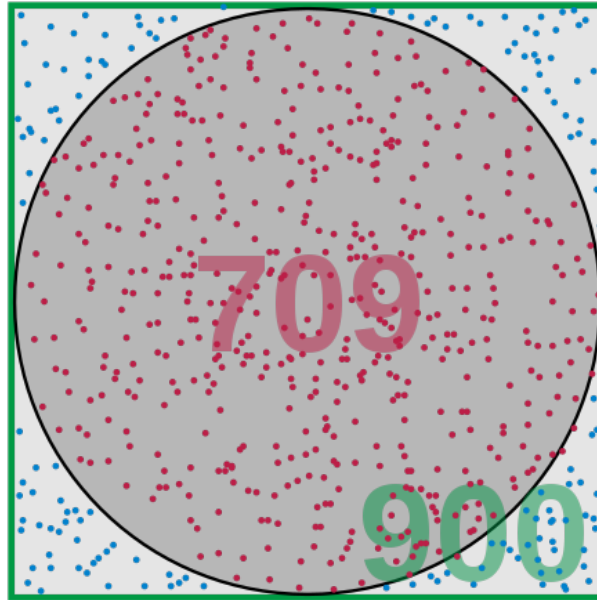


Figure 4.4: Illustration of the principle of Monte Carlo integration. If the area of the rectangular phase space is known, the area of the circle can be estimated from the ratio of randomly sampled points falling inside or outside the volume. In this example, the ratio of accepted points (709, red) to the total points sampled (900) approximates the analytic result for the ratio of the areas, $\pi/4$. Reproduced under public license from Ref. [153].

in the realistic environment of particle physics measurements. The MC integration approach is well-suited for this, as the procedure of randomly sampling the integrand over a phase space can be leveraged to draw events from the distribution. In this procedure, specific realizations of particle states, including quantum numbers and four momenta, are sampled from the possible outcomes with rate proportional to the probability for production, given by the phase space and squared amplitude of Equation 4.5.

4.5 Event simulation for experimental analysis

Calculations using the previously described techniques, which truncate at a specific order of perturbation theory and treat quarks and gluons as free particles, are referred to as “fixed-order” calculations. For the $pp \rightarrow WZ \rightarrow \ell\nu\ell'\ell'jj$ process, Equation 4.5 can be used to compute observables at fixed order for the underlying subprocesses, e.g., when a qq' interaction takes place and $jj = qq''$. This parton-level process can be used to derive useful predictions, however, events containing free partons and leptons do not offer a full picture of the objects that interact with the CMS detector. High-energy partons radiate quarks and gluons or split to $q\bar{q}$ pairs, reducing their energy until they reach the energy scale of hadronization Λ_{QCD} . Similarly, leptons can radiate photons, which can subsequently pair produce photons. Both effects are considered in the QCD and QED showers, which simulate successive soft emissions of massless or nearly-massless particles. At the scale Λ_{QCD} , the free partons form hadrons. Long lived hadrons interact with the detector, whereas shorter-lived hadrons decay to neutrinos, leptons, or other hadrons before interacting. These considerations, combined with additional pp interactions in a single LHC bunch crossing, as well as multiple parton-parton interactions in a single pp collision, contribute to the full $\mathcal{O}(1000)$ particle multiplicity measured in the CMS detector per event. The properties and kinematic distributions, as well as their interactions with the CMS detector, must be well-understood for a full simulation of a realistic event.

4.5.1 The parton shower and matching to matrix element calculations

The shower allows radiations at low transverse momentum and angular separation to be summed to all orders in perturbation theory. It is based on the observation that

the matrix elements $\mathcal{M}_{a \rightarrow b}$ and $\mathcal{M}_{a \rightarrow bc}$ are related as $\mathcal{M}_{a \rightarrow bc} \approx f(P_{b \rightarrow bc})\mathcal{M}_{a \rightarrow b}$ in the limit where p_T^b and the angular separation θ_{bc} is small [18]. In this expression, $f(P_{b \rightarrow bc})$ is an exponentially decreasing function, strictly less than one, of the splitting function $P_{b \rightarrow bc}$ for the particle b to bc . Thus, in this soft and collinear limit, additional radiation can be interpreted as a matrix element for production of the state b , and a factorizable probability for the incoming partons and final state particles—quarks, gluons, leptons, and photons—to radiate or pair convert. The MC techniques described previously are well suited for event-wise simulation of the shower, because random numbers can be linked to probabilistic parton splitting.

The shower is most important for QCD radiation, because gluon radiations and splittings occur with high probability. Therefore, it is generally referred to as a “parton shower,” even though photon radiation is also considered. The QCD splitting functions $P_{b \rightarrow bc}$ are equivalent to the Altarelli–Parisi functions used in the determination of the PDF, thus, they are well established. The process is followed iteratively, allowing subsequent radiations to evolve the state a into a state $a + X$, where X represents additional quarks or gluons. In this way, the parton shower evolves inclusive matrix element calculations, which give predictions for *at least* n partons ($n = 2$ for WZjj production), to be evolved to exclusive predictions for *exactly* n partons or clustered jets.

If the matrix element calculation includes outgoing partons, care must be taken to avoid overlap between the matrix element and shower components of the event. At LO, this is trivially accomplished by limiting the p_T of partons generated by shower splittings to be less than those generated by the matrix element—otherwise the harder shower emission would duplicate a possible configuration with harder matrix element partons and soft shower radiation, and the contribution would be double counted. This procedure, referred to as shower matching, is more complex for NLO calcula-

tions, because shower radiation from the Born diagrams (Fig. 4.3, left) overlaps the contribution of radiation from the real-emission diagram (Fig. 4.3, center). Two techniques were developed to resolve this conflict and enable showered and hadronized NLO predictions: the POWHEG [154, 155] and the MC@NLO [156] approaches.

The MC@NLO approach uses an analytic calculation of the parton shower to correct the matrix element predictions per event. This expected contribution is subtracted from the matrix element events, modifying the kinematics. The stochastic parton shower is then applied to the corrected events, which are not physically meaningful without the shower. The combined matrix element plus shower events, however, are accurate to NLO in QCD. The POWHEG technique takes a converse approach. It corrects the matrix element events to account for the parton shower-like behavior at low energy. The parton shower is applied to the corrected events, but any shower radiation harder than the real-emission matrix element is vetoed in favor of the matrix element contribution. We study both techniques for the $pp \rightarrow WZ$ process, and find agreement for inclusive variables.

In addition to more accurate cross section predictions, an NLO calculation has the advantage of describing additional parton contributions to the state with matrix-element-level calculations, via the real emission diagrams. Because calculating the virtual contribution contributing at NLO is a major challenge, it is natural to consider a way to combine the Born and real-emission contributions without relying on difficult loop diagram calculations. However, it is important to reiterate that matrix element calculations at a fixed order in perturbation theory are inclusive quantities. Tree-level diagrams that contribute at a higher order of α_s cannot be naively calculated and added to the prediction, without including the virtual contribution, that effectively regulates the relative contributions of the two. The parton shower prescription allows the inclusive prediction to be evolved into an exclusive one, with a prescription to

determine if a state remained a n -parton state or evolved to an $n + 1$ -parton state. Techniques known as matrix element merging exploit the parton shower paradigm to combine tree-level contributions with additional partons. In this approach, the parton shower determines how the inclusive state evolves from an inclusive to exclusive state of many partons, and the tree-level matrix elements are used to correct the distributions of the multi-parton states. This idea has also been extended to NLO calculations, avoiding the double virtual and real-virtual contributions in the FxFx [157] and MINLO [158] algorithms.

We use LO merging, via the MLM merging algorithm [159] implemented in MADGRAPH5_aMC@NLO, to simulate the QCD WZ process. By merging contributions to WZ production with up to three outgoing partons, an accurate description of the two-jet kinematics of the WZjj state is achieved without relying on computationally prohibitive higher-order calculations. We also exploit NLO-merging, but the number of outgoing partons that can be simulated is lower. The FxFx algorithm is used for alternative simulations of the QCD WZ process, for background diboson process, and for some backgrounds with top quark production. These techniques have been extensively validated with measurements of W and Z production [160], as well as in diboson states at in pp and $p\bar{p}$ collisions [161, 162].

4.5.2 The underlying event, hadronization, and decay

Hadronization refers to the transition from colored partonic states to the colorless hadronic states that populate the low-energy physical world. In general, the distributions of the total momentum of hadrons forming jets is approximately described by the momentum distribution from perturbative calculations of partons, so fixed-order calculations are still useful for predictions about the WZjj state. Generic parton-to-hadron fragmentation functions are also sometimes used to estimate the transition

from partons to hadronic jets to improve on these calculations. However, a much richer picture of hadron formation is obtained by simulating hadronization from the multi-parton state of the remnants of the interacting protons and the parton shower evolved to a cutoff scale $\sim \Lambda_{\text{QCD}}$ [163].

Partons do not hadronize individually, rather, they interplay to build colorless systems. To good approximation, the hadronization process can be considered universal, that is, the formation of hadrons is independent of the production mechanism of the partonic state, only on its configuration at the hadronization scale. Two models are used in general-purpose MC generators: the cluster model, implemented in SHERPA and HERWIG, and the Lund string model, implemented in PYTHIA. The cluster model is based on forming clusters of colorless quarks at the shower cutoff scale, after splitting gluons to $q\bar{q}$. As demonstrated by Amati and Veneziano in 1979 [164], color singlets at the cutoff scale have a mass distribution independent of the energy scale of the production process. The color singlet clusters are then decayed into hadrons according to the available phase space. The string model seeks to model the process of confinement more directly. The $q\bar{q}$ pairs are considered connected by confinement strings with potential energy proportional to the separation in the quark rest frame, $V = \kappa z$. Observationally, $\kappa \approx 1 \text{ GeV/fm}$, which is consistent with calculations in lattice QCD. As the $q\bar{q}$ pairs are separated, the energy increases, and it may be energetically favorable to snap the string with new $q'\bar{q}'$ pairs from the vacuum, as illustrated in Fig. 4.5. High energy gluons show up as kinks in the string, which give rise to more hadron formation in the gluon direction. The process continues until the color singlet pairs, or colorless qqq groups, are not energetic enough to break the confining string and form hadrons.

The cluster and string models both require input parameters derived from data, especially from measurements of jet formation in e^+e^- collisions at LEP. The string

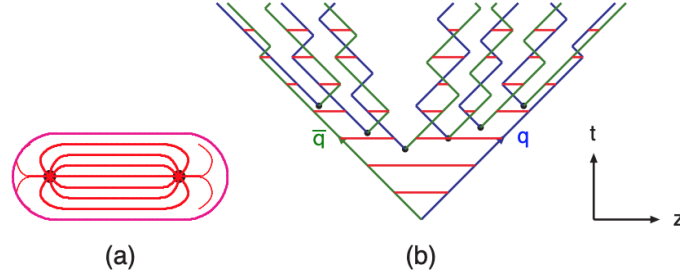


Figure 4.5: Illustration of hadronization in the string model. Field lines from gluon exchange lead to a string-like linear force between $q\bar{q}$ pairs (a). As the energy of the string increases, it is energetically favorable to produce new $q'\bar{q}'$ pairs (b).

model of PYTHIA is often seen to perform slightly better, however, it is dependent on a greater number of parameters derived from data. Good agreement is seen between these models for the WZjj state considered in this analysis. The simulation of the parton shower and matrix element merging, discussed in the previous section, have more impact on the dijet kinematics of WZjj events, which are critical to the analyses considered in this thesis, than does hadronization process.

Multi-parton interactions (MPI), remnants of the interacting protons, and pileup events from additional pp interactions are important for a realistic picture of particle formation in LHC collisions. The MPI is primarily soft; it does not generally result in easily identifiable jets, but it impacts the total energy transfer in a scatter event and affects the color connections of the scattered partons and proton remnants. The MPI and the remnants of the proton after the hard interaction impact the hadronization of the event; together they are referred to as the “underlying event.” For simulations in this thesis, the underlying event is modeled in PYTHIA, with parameters tuned by the CMS experiment using 7 and 8 TeV data as presented in Ref. [165], referred to as the CUEP8M1 tune, or with the default underlying event tunes of HERWIG or SHERPA. Pileup is simulated using large datasets of events of elastic and low-energy transfer pp scattering events, referred to as minimum bias events, simulated with PYTHIA 8.

Events are sampled randomly from the dataset and mixed into the simulated sample of hard-scattering events such that the number of pileup events matches that observed in data.

Hadron decays are simulated with a combination of theoretical models, including simple matrix element calculations for the light mesons and tau, and observational data from experimental measurements as tabulated by the Particle Data Group in Ref. [49]. Spin-dependent decays depend on the production mechanism of the particle. The production and decay spin correlations are taken into account by using MADSPIN [166] together with PYTHIA for simulations used in this analysis.

4.5.3 Detector simulation

To build a dataset of realistic collision events, used to guide this analysis and model expected signal and background distributions, the interactions of the particles that reach the CMS detector in the active and inactive volume are simulated using a detailed model in the GEANT4 program [127, 128]. GEANT4 includes a large set of physics models over an energy range from several eV to TeV, and it effectively serves as a repository of the known set of interactions of particles with matter. The paths of primary and secondary particles through the CMS detector are traced by GEANT4, using the MC method to determine interactions according to their probabilities. Decays of semi-stable particles are calculated and their decay products are tracked. Energy deposits and hits in the detector are recorded and transmitted to customized digitization software. The digitized information, designed to accurately emulate the real CMS electronics, is fed through the same reconstruction algorithms as real data.

The performance of the detector modeling and digitization have been tuned and validated using both test beam [167] and collision data [168]. Nonetheless, differences

in the simulated and collected data are unavoidable. They arise due to, e.g., inactive regions in the detector, shifts in the detector alignment, and radiation damage. Differences in data and simulation efficiencies are corrected using known quantities, such as the lepton kinematics from Z decays or jet energies in dijet events. These corrections are generally at the percent level. As discussed in Chapters 6 and 7, the agreement between data and MC simulation is excellent for most processes considered in this result.

4.6 Simulations of WZjj production

The primary generators used in this analysis are MADGRAPH5_aMC@NLO and POWHEG 2.0 for simulation of the hard-scattering interaction at LO or NLO, and PYTHIA 8 for parton showering and hadronization. The specific samples used for background processes are discussed in Chapter 6. Because the MC simulations are fundamental to distinguishing the EW- and QCD-induced contributions to WZjj, particular care is taken in studying the simulation of these processes.

Standard techniques to evaluate the uncertainties of MC simulations involve varying the parameters of the PDF fit and the scale of the μ_F and μ_R parameters, as discussed in Chapter 6. It is well established that these approaches offer a probe of the precision of a perturbative calculation. However, the extent to which they represent an uncertainty that can be interpreted in a statistical sense, and that fully covers the impacts of additional missing orders of the calculation, are unknown. In addition, additional parameters in a full simulation, such as the tuning of the shower and hadronization to data, are not captured by this approach. Even for fixed-order calculations, several parameters must be input to the calculation, including the functional form of the μ_F and μ_R , the couplings, and particle masses and widths.

A useful approach to obtain a broad picture of the uncertainty of theoretical predictions for a process is to compare results between different calculations. In principle, all differences should be traceable to specific choices in the calculations. However, probing all possible sources of differences in a hadronized prediction is a very tedious process. It is essential when understanding sources of disagreement, but may not be necessary when demonstrating that broadly different models of, e.g., parton shower and hadronization, give similar results. For the work presented in this thesis, we use comparisons of MC generators to demonstrate that the features of the EW WZ and QCD WZ processes that we exploit for this analysis are theoretically well understood. Much of this work was performed in collaboration with experts in the theoretical community at the Les Houches 2017 workshop, and reported in Ref. [169], in a study the author of this thesis proposed and co-convened.

Following this work, results are presented for the following event selection, referred to as the fiducial region.

- All charged leptons are required to have

$$p_{\text{T}}^{\ell} > 20 \text{ GeV}, \quad |y_{\ell}| < 2.5. \quad (4.6)$$

- For the leptons of opposite charge and same flavour, an invariant mass cut to single out the Z boson resonance is applied:

$$76 \text{ GeV} < m_{\ell^+\ell^-} < 106 \text{ GeV}. \quad (4.7)$$

- hadronic or partonic particles are cluster into jets using the anti- k_{T} algorithm [170] with radius parameter $R = 0.4$. At least two jets are required to have

$$p_{\text{T}}^{\text{j}} > 30 \text{ GeV}, \quad |y_{\text{j}}| < 4.7, \quad \Delta R_{\text{j}\ell} > 0.4, \quad (4.8)$$

and are called tagging jets.

- On the two leading tagging jets, typical VBS cuts are applied:

$$m_{jj} > 500 \text{ GeV}, \quad |\Delta\eta(j_1, j_2)| > 2.5. \quad (4.9)$$

4.6.1 Predictions from fixed-order calculations

Calculations of the EW WZ process at NLO in QCD were first performed more than ten years ago [149], however, only recently have these results been implemented in a form that is compatible with simulations of the parton shower and hadronization [171]. Recent results considering both NLO EW and QCD effects for WZjj production are also limited to fixed order [172]. Therefore, it is necessary to consider fixed-order results in order to study the best available predictions for this process. Moreover, the additional complications of showered and hadronized events can complicate efforts to establish the relative agreement of generators.

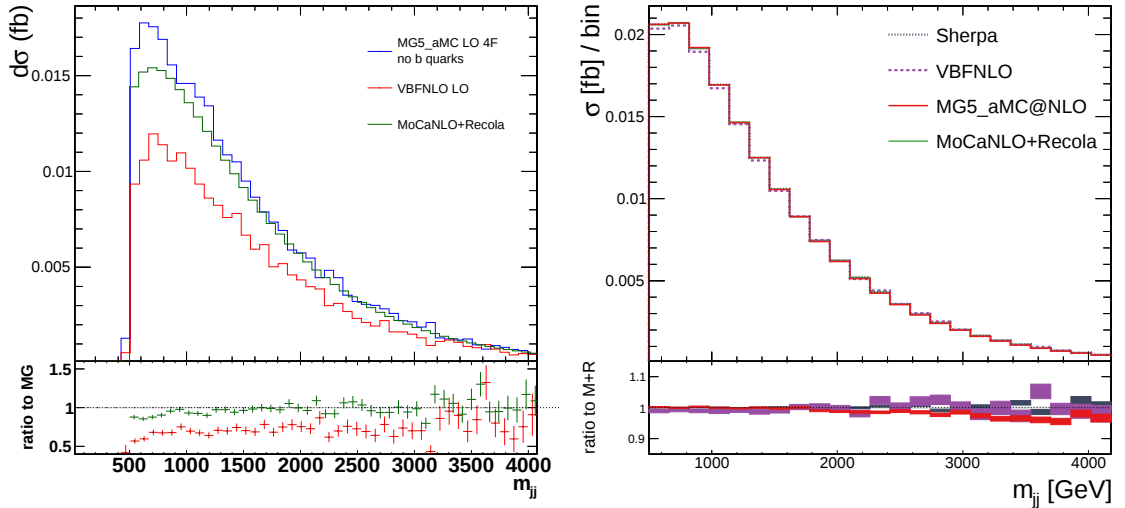


Figure 4.6: The fixed-order prediction for m_{jj} from MADGRAPH5.aMC@NLO, SHERPA, VBFNLO, and MoCaNLO+RECOLA at LO for the EW WZ process. The left comparison is made using the default settings of all input parameters for the public MC programs. The right plot is shown with parameters tuned to the values given in Ref. [169].

For this reason, we first consider comparisons of the EW WZ process at fixed order. Because the process has no α_s couplings at LO, the calculation has almost no μ_R dependence. The μ_F dependence is also relatively low, $\mathcal{O}(10\%)$ leading order. Furthermore, the NLO corrections to the cross section are small, the k -factor $k \equiv \sigma_{NLO}/\sigma_{LO} = 0.98$ in the fiducial region defined in Equations 4.6-4.9, calculated with VBFNLO [173]. Nonetheless, the process has a significant dependence on the masses, widths, and couplings of the particles used in the calculation, as demonstrated in Fig. 4.6. The Les Houches study demonstrated that discrepancies may arise at fixed-order when these parameters are not controlled, however, when using a common, well-motivated parameters (following the best-available measurements), excellent agreement is obtained. Because most settings can be determined in an objective way, we do not consider such differences to constitute additional uncertainty in modeling this process.

4.6.2 Hadronized predictions and fully-simulated events

We also study the modeling of the EW WZ process using hadronized events. The results are largely consistent with the studies at fixed order; agreement is seen across a broad range of generators for leptonic variables, or kinematic variables of the two-jet system. Specifically, the variables m_{jj} , $\Delta\eta_{jj}$, and $\eta_{3\ell}^* \equiv \eta^{3\ell} - \frac{1}{2}(\eta^{j_1} + \eta^{j_2})$ (sometimes referred to as the "Zeppenfeld variable") are all well described, and can be exploited in this analysis with limited uncertainty. Variables depending on more than three clustered jets, produced purely by the shower and hadronization for the MCs we consider, can have significant differences. This is expected: the parton shower and hadronization of SHERPA, HERWIG, MADGRAPH5_aMC@NLO, and PYTHIA use different models and different tuned parameters. Because many are phenomenological, it is difficult to constrain them extensively. We acknowledge a higher modeling uncer-

tainty for such variables, and seek to avoid relying on their simulation in this analysis. Example distributions with large and small shower dependence are shown in Fig. 4.7.

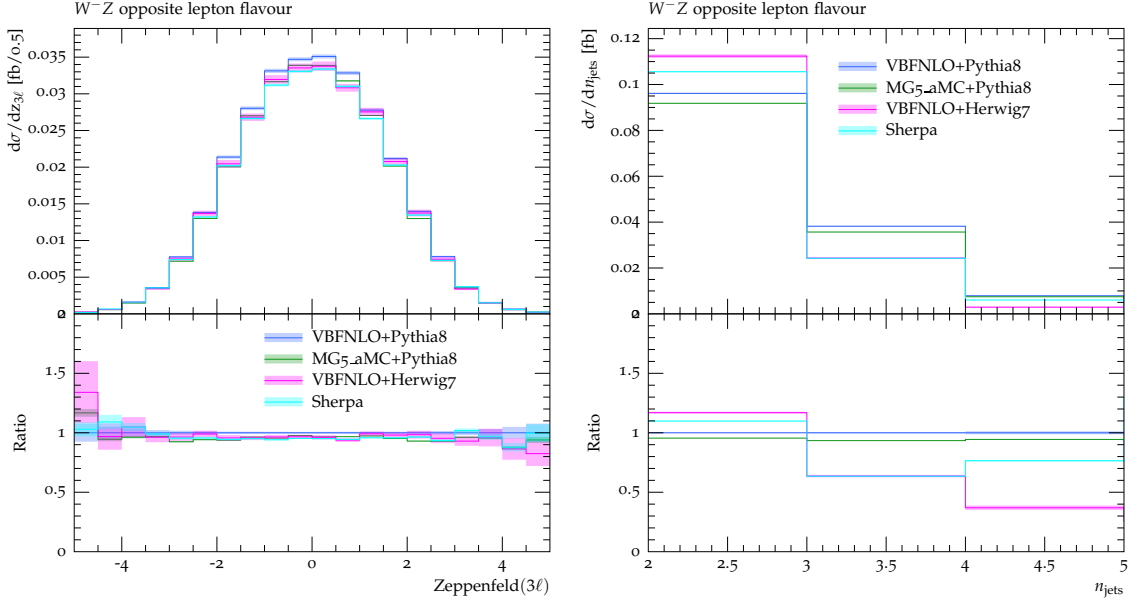


Figure 4.7: Predictions for EW $W^-Z \rightarrow e^+e^-\mu^-\bar{\nu}_\mu$, production, including the parton shower and hadronization. Variables well described by the fixed-order calculation, including $\eta_{3\ell}^*$ (left), are not significantly impacted by the shower and hadronization. Variables with a large dependence on shower and hadronization, including the number of jets with $p_T > 30$ GeV (right), show larger variance between MC generators.

The MC simulation used to guide this analysis and to extract results is generated with MADGRAPH5_aMC@NLO v2.4.2. We use the default dynamic scale choice and input parameters of MADGRAPH5_aMC@NLO. This choice was made for practical reasons, though we validate that only the total cross section, but not sensitive distributions, is affected by this choice, as illustrated in Fig. 4.8. The cross section with this configuration is about 10% lower than with the Les Houches settings. The EW WZ cross section measurement relies on shape predictions, but not on the predicted cross section, so this difference is not considered as an uncertainty. The shapes of two-jet variables are only majorly impacted by the corrections, is shown in Fig. 4.9 for EW WZ production at NLO with parton shower and hadronization from POWHEG [171]

These comparisons give us confidence that differential predictions for EW WZ production are well understood. The k -factor in this region is slightly lower than the VBFNLO fixed-order prediction.

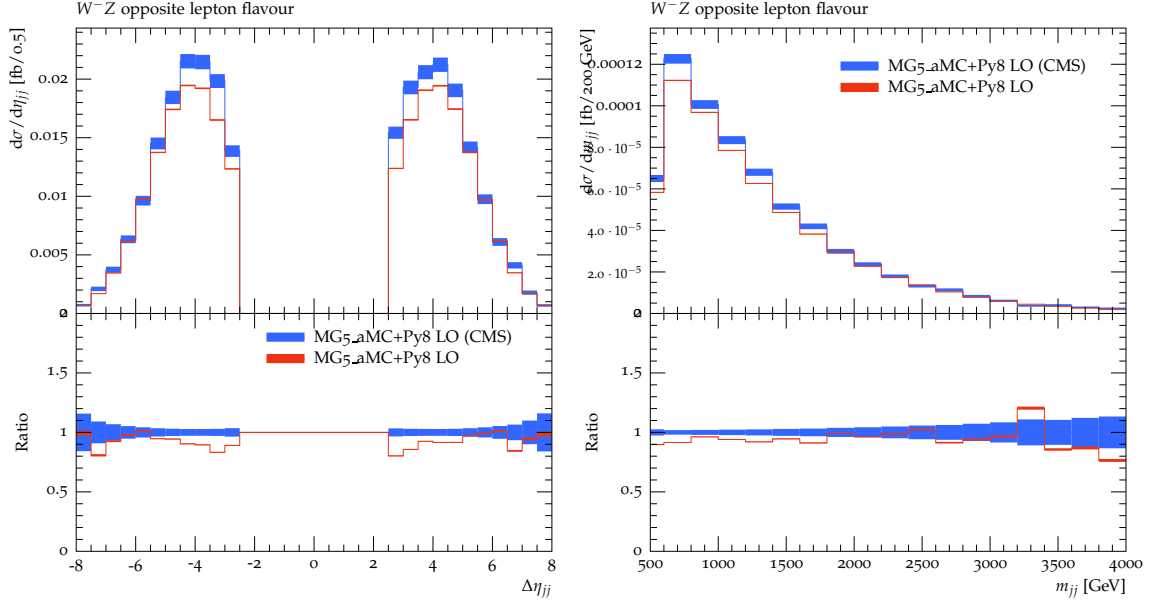


Figure 4.8: Predictions for EW $W^-Z \rightarrow e^+e^-\mu^-\nu_\mu$ production at LO in QCD simulated with MADGRAPH5_aMC@NLO v2.4.2 for the fiducial selection described in the text. The simulation used to extract results for this analysis is shown in blue. The simulation configured with the suggestions of the Les Houches study is shown in red.

Because QCD-induced WZjj production is $\mathcal{O}(\alpha_s^2)$ it has a large dependence on μ_F and μ_R , especially at LO. Fixed-order calculations have been performed at NLO [174], but hadronized simulations are computational challenging and are therefore not considered here. Several MC simulations of the QCD WZ process, including parton shower and hadronization, are considered. The simulations are inclusive in the number of jets associated with the leptonically decaying W and Z bosons, and therefore include the WZjj contribution studied in this analysis. The primary MC sample is simulated at LO with MADGRAPH5_aMC@NLO v2.4.2, with contributions to WZ production with up to three outgoing partons included in the matrix element calculation. The different jet multiplicities are merged using the MLM scheme [159]. An

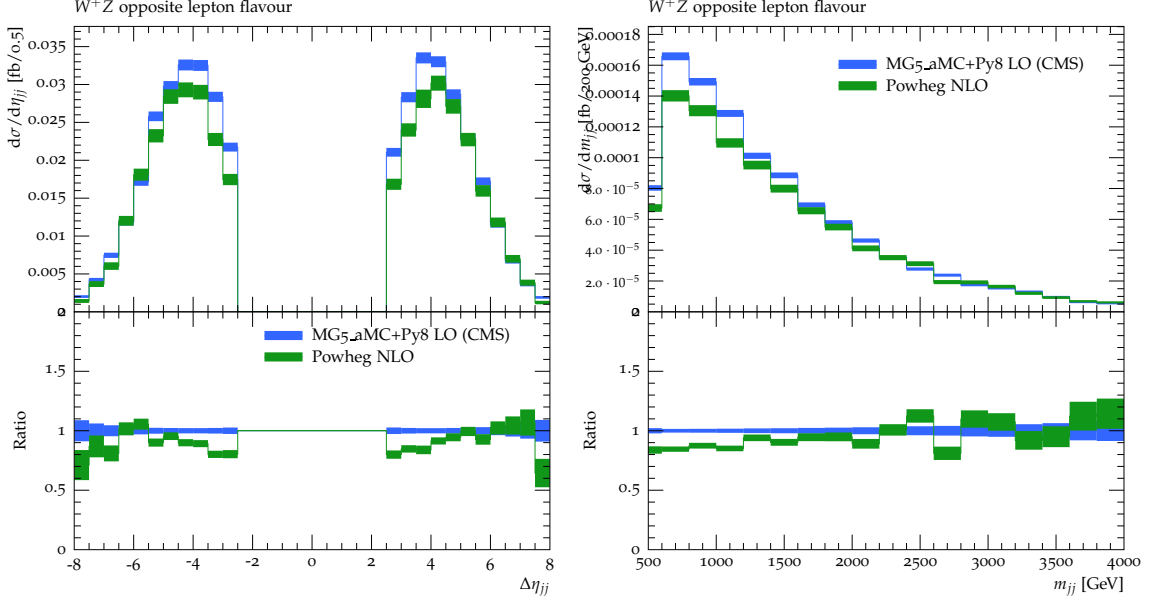


Figure 4.9: Predictions for EW $W^-Z \rightarrow e^+e^-\mu^-\nu_\mu$ production simulated at LO with MADGRAPH5_aMC@NLO v2.4.2 and at NLO with POWHEG for the fiducial selection described in the text.

NLO sample from MADGRAPH5_aMC@NLO v2.3.3 with zero or one outgoing partons at Born level, merged using the FxFX scheme, and an inclusive NLO sample from POWHEG 2.0 [154, 155, 175, 176] are also considered. The LO MC with MLM merging, referred to as the MLM-merged sample, is taken as the central prediction due its inclusion of WZ plus three-parton contributions at tree level, which are relevant to WZjj production. The other samples, which are used to access the modeling uncertainty in the QCD WZ process, are referred to as the FxFX-merged and the POWHEG samples, respectively.

The scale choice and input parameters are set to the default values of the respective generators. Due to the large scale dependence, the predicted cross sections for the WZjj state differ by up to 20% across the three MC generators considered. However, as for EW WZ, we are most concerned with the differential predictions for the process, because the total rate of WZ production can be constrained by the

data. Predictions for these MC generators, normalized to the prediction of MADGRAPH5_aMC@NLO in after the two-jet selection, are shown for variables sensitive to the EW WZ process in Fig. 4.10, where some differences are expected due to the parton shower schemes. Disagreement between predictions is largely within the theoretical uncertainties considered. The estimation of these uncertainties, and the additional uncertainty we derive from these comparisons after exploiting the data to constrain the predictions, are discussed in Chapter 6.

These studies of the MC simulations considered in this work have established that reliable predictions and uncertainties are observed for several variables sensitive to distinguishing the EW WZ and QCD WZ processes. In particular, we exploit the $\Delta\eta_{jj}$, m_{jj} , and $\eta^{3\ell} - \frac{1}{2}(\eta^{j1} + \eta^{j2})$ in the extraction of results using the central predictions of MADGRAPH5_aMC@NLO. The uncertainty in these predictions is in the range of 10–20% for the QCD WZ process, modeled with MADGRAPH5_aMC@NLO at LO with MLM merging, and 5–10% for the EW WZ process modeled with MADGRAPH5_aMC@NLO at LO. The uncertainties are discussed in more detail in Chapter 6.

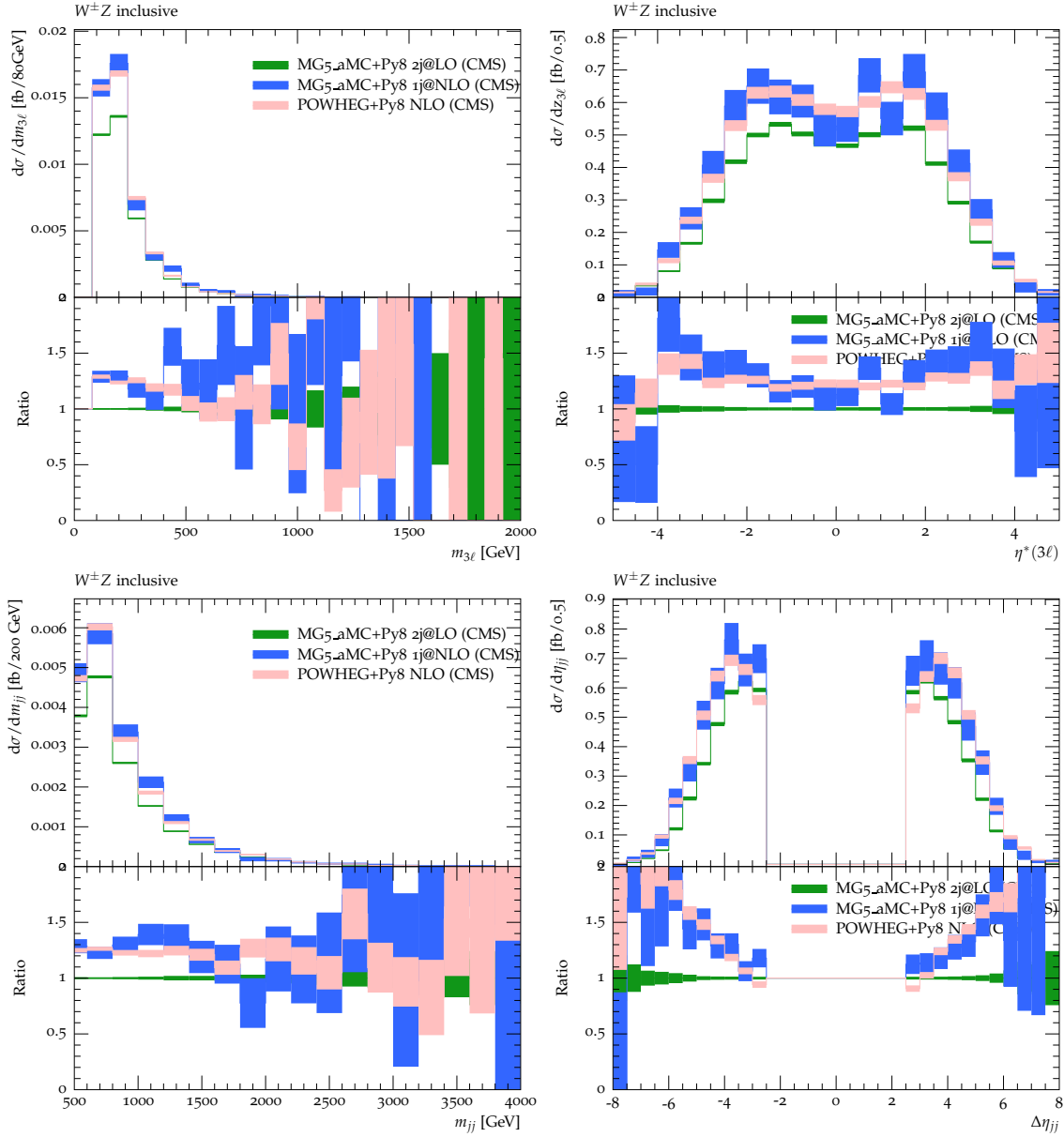


Figure 4.10: Comparison of predictions for the QCD WZ process using the POWHEG, MADGRAPH5_aMC@NLO MLM- and FxFX-merged samples in the fiducial region described in the text. The samples are normalized to the cross section from MADGRAPH5_aMC@NLO after selecting WZjj events.

Chapter 5

Event reconstruction

As described in Chapter 3, the CMS detector is built from complementary detector components with sensitivity to stable particles: electrons, muons, photons, and charged and neutral hadrons. Each can be distinguished by their likelihood of interaction and interaction characteristics with the detector components: all charged particles leave tracks in the inner tracker, charged hadrons deposit energy in the calorimeters, photons and electrons primarily deposit energy in the ECAL, and neutral hadrons interact primarily with the HCAL. The reality is more subtle than this idealized picture, illustrated in Fig. 5.1, but this forms the guiding principles in the *reconstruction* of particles from their interactions in the detector. After stable particles have been reconstructed, unstable particles are inferred from the stable particles to which they decay.

When an event is accepted by the HLT, the raw, digitized data from each detector subsystem is stored and simultaneously forwarded to a commercial processing farm for further processing. Detector-specific attributes, such as signal amplitudes, are converted to physical properties, such as energies, using calibrations established from controlled, independent measurements. System-wide patterns are built, from which

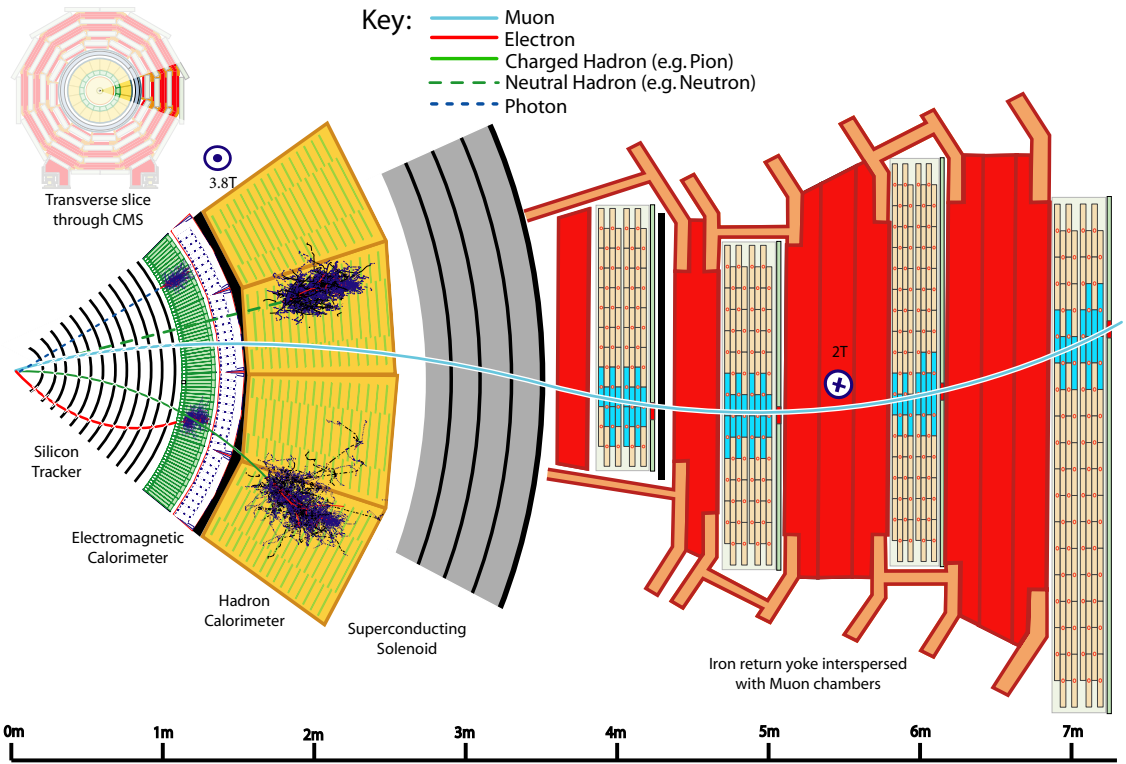


Figure 5.1: Idealized picture of particle interactions with the CMS detector subsystems.

a global picture of the event and its component particles is extracted. The reconstruction algorithms employed by the CMS Collaboration are built from work that began well before the LHC began taking data, benefiting from generations of previous experiments. The development of algorithms is an iterative process, with continual optimizations and incremental changes being made. The techniques described here relate to the data collection and reconstruction used by the CMS Collaboration in 2016. Many aspects of the reconstruction are shared by all analyses performed by the CMS Collaboration.

A degree of ambiguity in particle identification is unavoidable. Reconstruction *efficiency*, or a high rate of positive identification and reconstruction of real particles, must be balanced alongside the rate of misreconstruction and misidentification.

For example, electrons with a poorly reconstructed track will mimic photons, but because the track reconstruction efficiency cannot reach 100%, accepting only very high-quality tracks will lead to the rejection of real electrons if an interaction in the tracker is not readout. The optimal balance between efficiency and misidentification is dependent on the characteristics of the signal and background targeted by a particular analysis. In addition to the common reconstruction features, the identification criteria used in this work are also discussed.

5.1 Global event description via particle flow

The general characteristics of physics object identification outlined above are ubiquitous amongst past, current, and future collider detector reconstruction algorithms. Rather than applying these criteria object by object to select particle candidates of interest, the CMS Collaboration implements an ambitious *particle flow* (PF) algorithm that seeks to reconstruct and identify all objects in an event as particle candidates. This approach relies on a high-granularity detector and accurate tracking to connect features of different detector subsystem, which are combined for more accurate energy, momentum, and spacial measurements.

5.1.1 Tracks and vertices

Only energy deposits from charged particles traversing the pixel detector and strip tracker that exceed a threshold energy (adjustable for each element) are retained for offline storage. The position of each *hit* is locally reconstructed from clusters of adjacent pixel or strip elements based on the relative collected charge in the cluster. The efficiency for hits to be reconstructed is over 99% for active channels in both the pixel and strip detectors; approximately 98% of pixel channels and 96% of strip

channels were active during 2016 data collection. The hit position is calculated in local coordinates of the detector elements before being converted into a global position with respect to the full CMS detector. The transformation to global coordinates depends on the relative alignment of the individual tracker elements and their alignment with respect to the other detector systems. This is established through in situ measurements of the paths of cosmic rays to an uncertainty negligible with respect to the intrinsic position resolution of the silicon detectors [177].

The paths of charged particles through the detector are derived by fitting curves to the hits along a possible trajectory. The high combinatorics of possible track patterns that can be derived from the large number of hits arising from high-particle-multiplicity events is a computational challenge. To reduce the complexity of pattern identification while maintaining a high efficiency of reconstruction, the tracking finding algorithm proceeds in steps, designed to first identify unambiguous tracks before recovering those of lower quality. Hits used to form tracks are removed from consideration in the following steps, reducing the complexity and chance for misassociation.

In all steps, the algorithm begins with an initial projection of the track direction, referred to as the track *seed*. The track is then formed sequentially by extrapolating outward from the beam interaction—in the case of tracks seeded by the inner tracker, or inward from seeds in the muon chambers of calorimeter clusters. Subsequent hits are located by predicting their location with using the Combinatorial Track Finder (CTF) algorithm [119], an adaptation of the combinatorial Kalman filter algorithm [178]. At each step, the projected trajectory is updated to account for the additional associated seeds, simultaneously building and fitting the particle path. When the full trajectory through the tracker is established, the track is refit and discarded if it fails to meet quality criteria, such as an excessively high χ^2 of the fit.

The first track-forming iterations target high-quality, high- p_T tracks associated with the beam crossing region, or *beam spot*, established over many collisions by fitting the expected Gaussian beam profile to the observed pixel tracks. These initial iterations use seeds formed from hits in each layer of the pixel detector, relaxed to two hits in subsequent iterations. Following iterations relax the requirement for the track seed to be compatible with the beam spot, instead seeding via successive hits in the strip tracker, increasing sensitivity to displaced tracks from heavy-flavor hadron decays. Final iterations are seeded by the muon detectors using an outside-in approach.

Electrons traveling through the tracker have up to $\approx 85\%$ probability of radiating photons (known as bremsstrahlung) due to interactions with the significant tracker material. Therefore, electron tracks are likely to exhibit sharp kinks from sudden energy loss by bremsstrahlung. A modified algorithm known as Gaussian sum filter (GSF) [179] is thus employed to more accurately reconstruct electron-like tracks. In the CTF algorithm, the Kalman filter accounts for the material interactions with a Gaussian smearing around the expected hit location. In the GSF algorithm, a sum of Gaussian terms, designed to reproduce the Bethe formula for electron energy loss [180], is used. The GSF algorithm is applied to a subset of tracks formed via the CTF algorithm that have few associated hits or a relatively poor χ^2 to obtain an improved description under the electron hypothesis. The GSF algorithm is also used to form tracks seeded by energy clusters in the ECAL, described in Section 5.2.2.

Because the collisions analyzed for this work generally involve multiple pp interactions, tracks originate from a variety of vertices. Vertices associated with a pp collision and not an interaction or decay displaced from the collision point are identified by evaluating the common points of origin for a subset of the track collection that are well-reconstructed (low χ^2) and consistent with the beam spot. No addi-

tional condition on the track p_T is applied to select this reduced track collection. The deterministic annealing algorithm [181] is used to group the tracks into their most probably set of common vertices, before the vertex coordinates are derived by a fit to the associated tracks, considering their likelihood of misassignment to the vertex.

The *primary vertex* is the vertex associated with the hard scattering interaction of interest. The vertex that maximizes $\sum_i p_T^2$, where the sum considers the tracks associated with the vertex—clustered using the anti- k_T algorithm described in Section 5.2.3 with radius parameter $R = 0.4$ —and the transverse momentum imbalance associated with the vertex and clustered tracks. Studies of simulated WZ events demonstrate that this approach selects the correct pp vertex more than 99% of the time for signal processes considered here.

5.1.2 Muon system tracking

Charged particles traversing the muon system ionize the gas, leading to electrical signals on the wires and strips of the detector systems that are reconstructed as hits. As for the tracker, the hits are used to infer the paths and momenta of the interacting particles. All charged objects will be reconstructed by the system, but the large material presented by the tracker, ECAL, HCAL, and the support structure and magnet in front of the muon system ensures that the vast majority of the objects reaching the system are muons.

Hits in the individual DTs, CSCs, and RPCs are reconstructed using the principles described in Chapter 3. Straight-line tracks, known as “segments,” are built by fitting the hits across each chamber for the CSC and DT subsystems, which are built from multiple layers of independent gas volumes. The efficiency for the muon systems to reconstruct hits or segments from real muons traversing the system is evaluated using the “tag-and-probe” technique, for which the momentum and identification of

selected events can be inferred from independent characteristics to calibrate the detector performance. The hit efficiency of DTs is found to be greater than 97% over the entire system; for RPCs, it is greater than 94% (96%) in the barrel (endcap). The trigger requirement for the CSC chambers requires several consistent tracks, which biases efficiency measurements for individual hits. The efficiency to reconstruct segments through the chambers is instead considered. It is evaluated to be 97.4% efficient for real muons, determined from the percentage of “probe” muons reconstructed in the tracker that traverse the CSC detector to those with a reconstructed segment in the CSC chamber. Groups of DT and CSC segments are used to seed a Kalman-filter based tracking algorithm that builds “standalone” muon tracks from hits in the muon system [182].

5.1.3 Calorimeter clusters

The calorimeters measure both the location of an interacting electron, photon, or hadron, and its energy. Because energy is deposited in a shower with a significant depth and transverse extent, the energy must be collected across several crystals, scintillators, or fibers. An accurate energy measurement relies on collecting as close to 100% of the energy as possible. In addition, the distribution of the energy deposit in space can be used to identify the particle that induced the shower.

Clusters of energy are formed separately for the barrel and endcaps of the ECAL and HCAL and the ECAL preshower by combining the energy measurements in segmented components of the detectors. Reconstruction is seeded by detector elements with a measured energy deposit $E > E_{seed}$, from which *superclusters* are formed by aggregating neighboring detector cells that exceed energy threshold of twice the cell noise level with the seed cell. The E_{seed} values are optimized independently for the detector regions from simulation studies of photons and hadrons. The values range

from a few hundred MeV for the ECAL barrel and endcap up to over 1 GeV in the HCAL endcap; the exact values are given in Table 2 of Ref. [183]. Clusters are formed from the superclusters by assuming the energy deposit from an interaction is Gaussian distributed. The number of interactions contributing to the supercluster is obtained from a maximum likelihood fit that considers the number of Gaussian energy deposits and their amplitudes. In the case of multiple deposits, the distributions obtained from the fit are used to divide decompose the supercluster into individual clusters. Because of the rough segmentation in the forward region, HF clusters consist of a single segmented area.

5.1.4 Particle flow candidates

The principle role of the PF algorithm is to *link* the the inner tracks, muon system tracks, and calorimeter clusters to form *PF candidates*. The PF candidates are categorized as neutral or charged hadrons, electrons, muons, and photons that are used in to correct the energy response of the detector in a particle-specific way and to build the identified objects selected for analysis. The link algorithm evaluates to consistency of tracks and clusters by extrapolating from the last hit of each inner track to:

- The depth in the ECAL of the expected maximum energy deposit a typical electromagnetic shower
- The plane of the ECAL preshower layers
- One hadronic interaction length into the HCAL

Tracks and clusters are linked if a cluster is found within one ECAL or HCAL crystal. In the case of multiple matches, the cluster with its maximum closest to the track

position is taken. Clusters in the ECAL and HCAL are linked if the ECAL clusters overlaps the HCAL cluster. For GSF tracks linked to an ECAL cluster, energy deposits in the ECAL consistent with photon radiation tangent to the GSF at a silicon layer are collected and linked to the GSF track and ECAL cluster.

The linked tracks and energy clusters, referred to as *PF blocks*, are further analyzed to be cauterized as a specific particle candidate that is used in the final selection of events of interest. First, muons are identified by linked tracks in the tracker and muon system. Electrons and photons are then considered simultaneously. Electrons are seeded by GSF tracks linked to ECAL clusters whereas photons are seeded by ECAL clusters with no linked track. Candidates are classified as PF electrons based on a multivariate regression trained on characteristics of the GSF track and the energy cluster. The PF photons are categorized by isolation and the ratio of energy in linked ECAL and HCAL clusters.

For each step, PF blocks associated with PF candidates are removed before further categorization. After removing muons, electrons, and photons, remaining blocks are categorized as charged and neutral hadrons. Neutral hadrons are formed from clusters in the HCAL that are not linked to ECAL deposits, whereas charged hadrons require an ECAL and HCAL cluster link. The momentum of tracks linked to HCAL clusters is used to improve the poor resolution of the HCAL for a significant improvement below several hundred GeV. Because the most ubiquitous neutral hadron, π^0 , decays to two photons, neutral candidates are dominated by photons, which are measured with considerably better energy resolution than neutral hadrons. In the HF, the ratio of energy deposits in the long fibers to short fibers is used to categorize clusters as HF photon and charged hadron candidates. The PF identification is used to correct the energy assignment based on object-specific calibrations in the calorimeters.

5.2 Physics objects

5.2.1 Muons

Muons selected for this analysis must be identified as global muons by the PF algorithm described previously. They must be within the muon detector acceptance $|\eta| < 2.4$ and must have $p_T > 10$ GeV. Additional reconstruction criteria are applied to reject *nonprompt* muons, arising from hadron decays, and hadron *punch through* into the muon system, while maintaining a reconstruction efficiency greater than 96%. Following the *Tight* identification criteria of Ref. [184], with slight modifications, muon candidates must satisfy:

- The global muon fit must satisfy $\chi^2/\text{d.o.f.} < 10$ (where d.o.f is the number of degrees of freedom in the fit)
- The distance of the track vertex from the primary vertex in the x - y plane d_{xy} must satisfy $d_{xy} < 0.2$ mm
- The longitudinal distance (in the z plane) of the track vertex from the primary vertex plane d_z must satisfy $d_z < 1$ mm
- At least one hit in the pixel detector
- More than five hits in the strip tracker
- At least six hits in the tracker must be used in the track fit

To reject muons from decays of hadrons in jets, selected muons are required to be isolated from other PF objects in the event. The relative muon isolation, the ratio of the sum of the p_T in an η - ϕ region $\Delta R < 0.4$ around the muon is defined as

$$I_{\mu,rel} = \frac{1}{p_T^\mu} \sum_{\Delta R < 0.4} \left[p_T^{h^\pm, PV} - \max \left(p_T^{h^0} + p_T^\gamma - \frac{1}{2} p_T^{h^\pm, PU}, 0 \right) \right]. \quad (5.1)$$

Here $p_T^{h^{\pm,PV}}$ and $p_T^{h^0}$ are, respectively, the charged and neutral hadrons identified by the PF algorithm. The tracks of charged hadron candidates allow them to be associated to the primary vertex, $h^{\pm,PV}$ or to pileup interactions, $h^{\pm,PU}$. Such an assignment cannot be made reliably for neutral hadrons and photons, therefore, all neutral hadrons are included in the sum. To account for the expected pileup contribution, the sum is corrected by the term $p_T^{h^{\pm}}/2$, because particle production in pileup interactions is expected to be dominated by charged hadron production at nearly this ratio. A rough estimation of hadron production split evenly between π^+ , π^- and π^0 gives exactly this ratio, and detailed simulations confirm this expectation.

The $I_{\mu,rel}$ is used to define a baseline, or *loose muon ID*, and a signal, or *tight muon ID*. Explicitly, the criteria $I_{\mu,rel} < 0.4(0.15)$ is used for loose (tight) ID. Muons satisfying the loose ID, but not the tight ID, are used to estimate the residual contamination from nonprompt muons in the event selection. Tight muons are used to select signal-like events. A further description of the approach is given in Chapter 6.

5.2.2 Electrons

Electron identification for this analysis begins with the GSF tracks linked to the ECAL clusters by the PF algorithm. However, unlike for muons, selected electrons are not uniquely a subset of the PF electrons candidates. In particular, the BDT discriminant used by the PF algorithm to categorize electrons and photons is not considered. Instead, conditional criteria (referred to as “cuts”) are applied to specific characteristics of the electron candidate tracks and associated clusters.

Electron candidates must be within the ECAL detector acceptance $|\eta| < 2.5$ and must have $p_T > 10 \text{ GeV}$. As for muons, baseline (*loose*) and *tight* selections are defined. Electrons satisfying the baseline selection are used to estimate the residual contribution from nonprompt electrons, while electrons satisfying both the loose se-

lection and the full *tight* criteria are selected as signal candidate events in the analysis. The baseline selection was originally designed to mimic the electron selections applied at the HLT level for the tightest trigger paths used in this analysis, however, stronger conditions were applied at the HLT towards the end of data collection. Nonetheless, the efficiency loss is recovered by additional trigger paths as described in Chapter 6. The requirements placed on the selected *loose* and *tight* electron candidates are summarized in Table 5.1. The definitions and motivations of the variables shown are discussed in the following.

Conditions on the GSF track and linked PF superclusters in the ECAL and HCAL are applied to select candidates of interest. The track and cluster positions must be consistent, defined by the variables $|\Delta\eta_{seed}|$ and $|\Delta\phi_{in}|$. The $\Delta\eta_{seed}$ is the difference between the η direction specified by the track seed and the η of the PF cluster that seeds the ECAL supercluster, and $\Delta\phi_{in}$ is the difference between the track seed extrapolated to the supercluster and the supercluster ϕ . The energy threshold for electron detection ensures that the electron mass is negligible, so the momentum measured by the track curvature and energy measured by the ECAL should be nearly equal. This is ensured by requiring the variable $1/p_T^{track} - 1/E^{SC}$, where E^{SC} is the energy of the supercluster, be close to zero. Electrons deposit the majority of their energy in the ECAL. Therefore, the ratio of the hadronic to electromagnetic energy H/E , measured in terms of the PF ECAL and HCAL clusters, should be small. To ensure a well-reconstructed track, the baseline (loose) electron selection requires a good GSF track χ^2 . For the tight identification, the track is required to have at most one missing hit in the tracker, and the electron candidate is rejected if any tracks consistent with a photon-to- e^+e^- conversion are found.

The variable $\sigma_{in\eta}$ is used to assess the shape of the shower in the η . It is defined

as

$$\sigma_{i\eta i\eta} = \sqrt{\sum_{i \in 5 \times 5} \frac{w_i(i\eta_i - \bar{i\eta})}{w_i^2}}, \quad (5.2)$$

where the sum is over the 5×5 block of ECAL crystals centered around the cluster seed, and the weights w_i are defined as

$$w_i = \max\left(0, 4.7 + \ln \frac{E_i}{E_{5 \times 5}}\right). \quad (5.3)$$

The variable $i\eta_i$ is the η of each crystal in the sum, expressed in integer units numbering the crystals, a choice that allows the variable gaps between crystals to be ignored. In the endcaps, the arrangement of crystals is in the x - y plane, and $i\eta$ with crystal integer numbering ix . The $\bar{i\eta}$ is the $i\eta$ of the 5×5 cluster, E_i is the energy of the i th crystal, and $E_{5 \times 5}$ is the energy sum of the 5×5 cluster.

As for muons, selected electrons must be isolated from other event activity. The isolation is defined as

$$I^\ell = \left(\sum p_T^{\text{charged}} + \max\left[0, \sum p_T^{\text{neutral}} + \sum p_T^\gamma - p_T^{\text{PU}}\right] \right) / p_T^\ell, \quad (5.4)$$

where the sum considers objects contained in a cone of $\Delta R < 0.3$ around the electron. The exact definition depends on the physics objects considered in the sum, and the approach to pileup subtraction. For the loose electron selection, the isolation defined in terms of the tracks and in terms of the PF energy clusters in the ECAL and HCAL are considered separately. The track isolation considers only charged tracks in the tracker, neglecting the neutral component and neglecting corrections for the isolation. For the PF cluster isolations, the sum of ECAL and HCAL cluster energies are corrected by estimated the pileup energy contribution per event using the “jet area” technique proposed for jets in Ref. [185]. In this approach, $p_T^{\text{PU}} \equiv \rho A_{\text{eff}}$, where the average transverse momentum ρ is calculated from the median of the ratio of the jet transverse momentum to the jet area, p_T^j/A_j , for all pileup jets in the event. The

effective area A_{eff} is the geometric area of the isolation cone times an η -dependent correction factor that accounts for the residual dependence of the isolation on the pileup. The A_{eff} is calculated separately for the terms (e.g., neutral or charged) considered in the pileup subtraction. For the tight electron identification criteria, the PF objects are used. Charged PF candidates are required to be associated with the primary vertex. Because the neutral objects cannot be directly associated with a vertex, so the pileup correction considers only the neutral component.

Table 5.1: The identification criteria used to select electrons for the estimation of nonprompt backgrounds (loose ID) and to select signal events (tight ID). The variable of interest specified in the leftmost column must be smaller than the values indicated to the right. The conditions on the barrel region ($|\eta| < 1.479$) and the endcap ($1.479 < |\eta| < 2.5$) are optimized independently.

Variable	Loose ID		Tight ID	
	ECAL barrel	ECAL endcap	ECAL barrel	ECAL endcap
$ \Delta\eta_{\text{seed}} $	0.004	—	0.00308	0.0605
$ \Delta\phi_{\text{in}} $	0.020	—	0.0816	0.0394
$ E_{SC}^{-1} - p_{\text{track}}^{-1} $		0.013	0.0129	0.0129
$\sigma_{i\eta i\eta}$	0.011	0.031	0.00998	0.0298
H/E	0.060	0.065	0.0414	0.0641
I^{ECAL}	0.160	0.120		1
I^{HCAL}		0.120		1
I^{track}		0.08		1
$I_{e,\text{rel}}^{\text{PF,corr}}$		—	0.0588	0.0571
GSF track $\chi^2/\text{d.o.f.}$	—	3.0		—
Max missing tracker hits		—		1

5.2.3 Jets

As introduced in Chapter 2 and discussed in Chapter 4, partons produced in the hard interaction are observed in the detector as clusters of hadronized particles known as jets. Perturbative predictions rely on a clustering algorithm that is insensitive to the low-angle and low-energy behaviour of partons. A variety of such “infrared-safe”

algorithms have been developed and studied; an overview can be found in Ref. [186]. Because it is advantageous to use a common jet-clustering algorithm for experimental measurements and theoretical predictions, the anti- k_T algorithm [170] is used by both the ATLAS and CMS Collaboration.

The anti- k_T algorithm sequentially clusters objects into composite jets based on their separation in rapidity y and azimuthal angle ϕ . A weighted distance parameter is defined by

$$d_{ij} = \min(k_{Ti}^{-2}, k_{Tj}^{-2}) \frac{\Delta R_{ij}^2}{R^2}. \quad (5.5)$$

where R is the radius parameter, k_{Ti} and k_{Tj} index the transverse momentum of the objects considered in the comparison, and $\Delta R^2 = (y_i - y_j)^2 + (\phi_i - \phi_j)^2$. A parameter $d_{iB} = k_{Ti}^2$ is defined as the p_T of the i th object with respect to the beamline. The values of d_{ij} and d_{iB} are computed for all objects in the event. If a value d_{ij} is the minimum parameter, the i th and j th objects are merged into a composite objects by summing the four vectors. If d_{iB} is the minimum, the object is removed from further iterations. The procedure is repeated until all objects in the event are clustered.

The anti- k_T algorithm has several favorable theoretical and experimental properties: it is infrared safe, results in a jet that is circular in the y - ϕ plane for hard partons i and j with $\Delta R_{ij} > 2R$, and is computationally fast in the **FastJet** [187] implementation. A radius parameter of $R = 0.4$ is used for this analysis. This is the choice made by the majority of 13 TeV studies made by the ATLAS and CMS Collaborations, as it is found to be a good balance between minimizing the jet size to reduce contributions from pileup and maximizing the jet size to capture all hadrons associated with the fragmentation of hard partons.

Selected jets used in this analysis are defined by clustering the PF candidates reconstructed in Section 5.1.4. To reduce the contributions from pileup, charged hadrons unambiguously not associated to the primary vertex are removed from the

PF candidates collection before clustering [188] (referred to as charged-hadron subtracted, or CHS jets). Leptons are not removed from the subset of PF candidates considered, however, clustered jets with $\Delta R(j, \ell) < 0.4$ are removed from the jet collection. The contribution from neutral pileup is subtracted from the jet energy using the jet area method described for lepton isolation.

The energy scale and resolution of the jets is corrected in data and simulation to account for the complex detector response and inefficiencies. The detector response is nonlinear and dependent on the jet composition, which may not be well-modeled by MC simulations, so the energy corrections are derived and applied to the clustered jets rather than the individual PF constituents. The corrections are primarily derived from MC simulation, using ratios of “generator” jets formed by clustering all stable ($c\tau > 1$ m) hadronized particles in the MC simulation to jets of PF CHS jets in the reconstructed simulation. The ratios are used to correct the energy of reconstructed jets in the MC simulation and in data. Residual differences in data and MC simulations are corrected for by measuring the jet energies in data and MC for samples of Z+jet, γ +jet, and dijet events. For single jet events, the energy of the Z bosons (with decays to electrons or muons) or γ should be balanced with the momentum of the recoiling jet. For dijet events, a well-measured jet in the central part of the detector is balanced with a forward jet to derive energy corrections in the region beyond the tracker acceptance [189].

The uncertainty in the procedure is estimated by assessing the uncertainty in each component of correction, including comparisons of the residual corrections in data and MC simulation with different data samples and variations of the pileup subtracting estimation [190]. The total uncertainty and the uncertainty by correction component are shown in Fig. 5.2. As illustrated, the uncertainty increases at high η to 5-6% for $p_T = 30$ GeV jets. Because forward jets are a key component of this analysis,

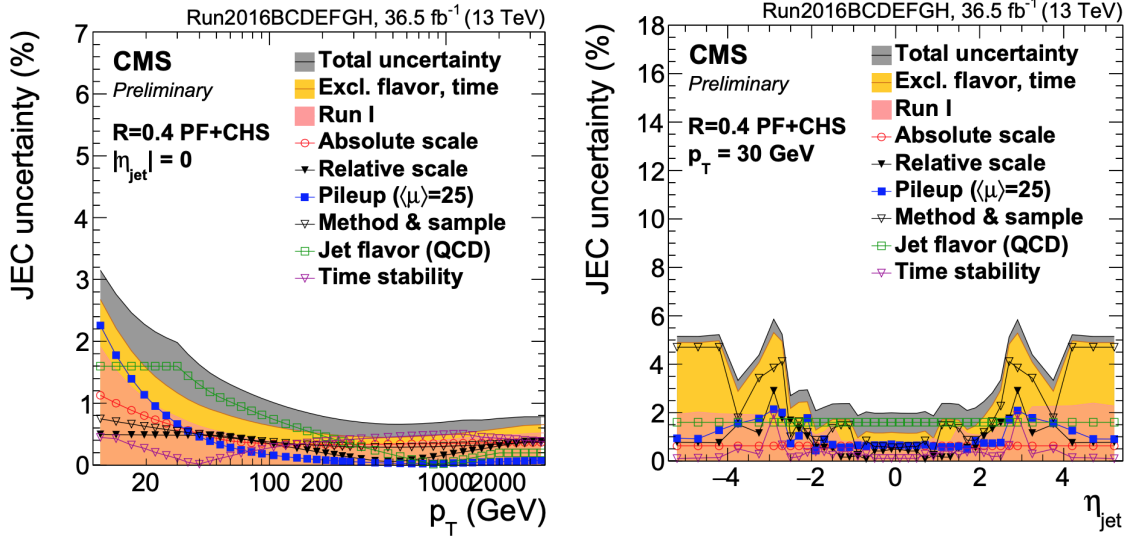


Figure 5.2: The uncertainty in the jet energy corrections (JEC) for PF charged-hadron-subtracted (CHS) jets vs. jet p_T (left) and jet η (right) for fixed jet $p_T = 30$ GeV. The total uncertainty and the uncertainty by component of the energy correction type are shown.

the jet energy scale and resolution uncertainties are the leading contributions to the measurement uncertainty.

5.2.4 Heavy-flavor jets

The hadronization of b quarks primarily produces baryons and mesons with larger masses and lifetimes than pions or other first- and second-generation hadrons. The decay lengths of such B-hadrons (e.g., B^0 , a $d\bar{b}$ bound state) produced at the LHC are typically on the order of centimeters [49]. The decay products of such B hadrons will be displaced from the primary vertex by this amount, which is within the resolution of the track reconstruction to be distinguished from the primary vertex. This observation is exploited to distinguish jets arising from the hadronization of b -quarks. Because the branching ratio $\mathcal{BR}(t \rightarrow Wb) \approx 1$, b -quark tagging is an effective means to identify top-quark production, a major background for this analysis.

The b quark jets are identified in this analysis via the combined secondary vertex b-tagging algorithm [191] with the *tight* working point defined in Ref. [192]. The algorithm using information on the PF candidate tracks and their associated vertices in the jet to train a multivariate discriminant, which is trained to distinguish b hadrons in MC simulation. The efficiency for selecting b quark jets with $p_T > 30 \text{ GeV}$ and $|\eta| < 2.4$ with the tight working point is $\approx 49\%$ with a misidentification probability of $\approx 4\%$ for c quark jets and $\approx 0.1\%$ for light-quark and gluon jets.

5.2.5 Missing transverse momentum

Neutrinos produced in the hard scattering interaction are inferred from the total imbalance in transverse momentum of the event, referred to as the *missing transverse momentum* $\vec{p}_T^{\text{miss}}$. It is reconstructed from the sum of all the PF candidates in the event,

$$\vec{p}_T^{\text{miss}} = - \sum_{\text{PF cands}} \vec{p}_{T,\text{cand}}. \quad (5.6)$$

Pileup collisions, which are generally soft hadronic interactions, do not produce significant $\vec{p}_T^{\text{miss}}$. Therefore, all PF candidates in the event are used in the sum, even those not associated with the primary vertex.

The measurement of the $\vec{p}_T^{\text{miss}}$ is improved by correcting jet energies with the procedure described in the previous section. The PF candidates that are clustered into jets are subtracted from the $\vec{p}_T^{\text{miss}}$ sum and replaced by the corrected energies,

$$\vec{p}_{T,\text{corr}}^{\text{miss}} = \vec{p}_T^{\text{miss}} - \left(p_T^j - \sum_{\text{cand} \in j} \vec{p}_{T,\text{cand}} \right). \quad (5.7)$$

The corrected $\vec{p}_T^{\text{miss}}$ is used in this analysis. Its uncertainty is dominated by the uncertainty in the jet energy corrections. A smaller uncertainty is associated with the *unclustered energy* of PF candidates not associated with a jet [193].

Spurious $\vec{p}_T^{\text{miss}}$ arises in events in which an object is significantly mismeasured, biasing the momentum sum of the event. Events with features of detector noise, such as large spikes in a single ECAL crystal or HCAL readout, as well as those with muon interactions in the calorimeters and CSC chambers that are symmetric in ϕ , indicative of *beam halo* muons produce in proton beam interactions with the beam pipe material, are vetoed.

Because the η direction of the $\vec{p}_T^{\text{miss}}$ vector is unknown, its magnitude p_T^{miss} is generally considered. In this analysis, the p_T^{miss} is used to select events with a leptonically decaying W boson giving rise to $p_T^{\text{miss}} \sim m_W/2$.

Chapter 6

Analysis Strategy

The EW WZ process exhibits a distinct experimental signature characterized by two forward, high-momentum jets, typically arising from the hadronization of two valence quarks from each of the colliding protons scattered away from the beam line, in addition to the vector bosons. We exploit the leptonic decay channel of the W and Z bosons in this analysis, as the three-lepton event signature is relatively free from other background processes. Events with exactly three leptons (electrons or muons), missing transverse momentum, and two jets at high pseudorapidity η with large dijet system invariant mass m_{jj} are selected for this analysis. A candidate EW WZ event selected by this analysis, demonstrating these characteristics, is shown in Fig. 6.1.

Deviation from the SM prediction in the rate of EW WZ production, or in kinematic distributions sensitive to the scattering energy of the process, such as m_{WZ} , could indicate the presence of BSM physics. This section describes the selection of events, the estimation of the expected signal and background contributions, the assessment of systematic uncertainties impacting these estimations, and the statistical techniques used to derive results. Results are presented in both SM and BSM interpretations, as described in Chapter 2.

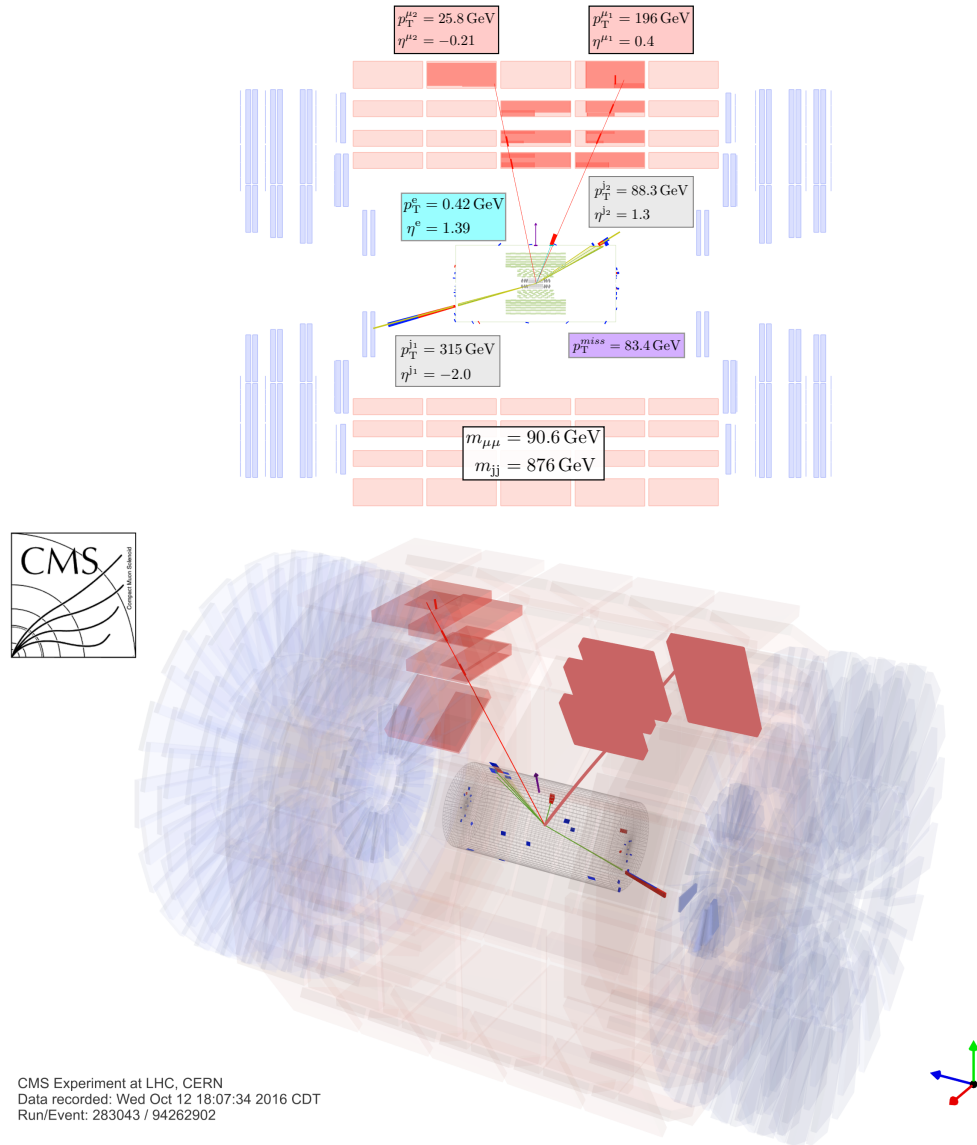


Figure 6.1: Visualization of a candidate EW WZ event selected by in this analysis. High momentum and well separated jets are observed in the forward part of the detector. Two muons (red) satisfy the p_T requirements while forming a composite object with invariant mass consistent with m_Z . Moderate missing transverse momentum (purple arrow) and an electron candidate (light blue) are consistent with a W-boson decay.

6.1 Background composition

Several process can mimic the $WZjj \rightarrow 3\ell jj$ signal. For the EW WZ selection described in Section 6.3, true WZjj production is estimated to account for 3/4 of the expected event yield, with background processes accounting for 1/4 of the expected events. The estimation of each background contribution is discussed in Sections 6.4.1 and 6.4.2.

We refer to processes with at least three leptons produced in the hard scattering interaction as “prompt.” This includes the leptonic decay channels of ZZ, tZq, $t\bar{t}Z$, and VVV production, where V refers to a W or W boson. The tZq process, where the top quark decays leptonically via a W boson, is the dominant prompt background contribution, accounting for 25% of background events after the EW WZ selection. In this process, one vector boson is radiated from the incoming quark, leading to a forward jet in addition to the three lepton signature. The ZZjj, $t\bar{t}Z$, and VVV processes enter the WZjj selection when leptons are outside of the detector acceptance, leading to the three lepton and p_T^{miss} signature. The $t\bar{t}Z$, ZZ, and ZZ processes are estimated to account for about 15%, 10%, and 5% of the total background respectively.

“Nonprompt” processes arise when leptons which are produced outside the hard interaction, especially through leptonic decays of hadronic states in jets and photon conversion to leptons in the detector material. These processes are dominated by Z+jet production, in which misidentification and mismeasurement of jet activity mimics a third lepton and gives rise to p_T^{miss} in the event, $t\bar{t}$ production, and $Z\gamma$ production. The expected contributions from these processes are 30%, 5%, and 10% of the total background for events satisfying the EW WZ signal selection.

The component of the WZjj process considered as signal or background is adjusted for the interpretation of results. For the WZjj fiducial cross section measurement, all SM WZjj production mechanisms are considered signal. For the EW WZ measure-

ment, the EW WZ process is considered signal, and the QCD WZ process is considered background. The QCD WZ process is the dominant background in this search, with a ratio of 3 : 1 to the EW WZ process in the signal region. For new physics searches, all SM WZjj processes are considered background. The ratio of EW WZ and QCD WZ processes is assumed to be consistent with the SM expectation.

6.2 Event triggering

Collision events are selected by triggers that require the presence of one or two electrons or muons. The p_T threshold for the single lepton trigger is 25 (20) GeV for the electron (muon) trigger. For the dilepton triggers, with the same or different flavors, the minimum p_T of the leading and subleading leptons are 17 (17) and 12 (8) GeV for electrons (muons), respectively. The combination of these trigger paths brings the trigger efficiency for three-lepton events satisfying the minimum p_T requirements to nearly 100%. However, partial mistiming of signals in the forward region of the electromagnetic calorimeter (ECAL) endcaps ($2.5 < |\eta| < 3.0$) led to a reduction in the level 1 trigger efficiency.

Larger-than-expected aging effects in ECAL crystals altered the energy pulse shape of signals in the ECAL, leading to early event readout for a significant fraction of events with high electromagnetic activity in the ECAL endcaps towards the end of 2016 data taking. This effect, known as pre-firing, is significant because requirements necessary to avoid excessive data readout prevent events from successive bunch crossings from being accepted by the CMS trigger system. That is, if an event has been accepted by the CMS trigger system, it is impossible for the following event to also be accepted.

In normal conditions, the inefficiency due to this requirement is very small and

unbiased. Systematic pre-firing, however, can lead to a biased and unrecoverable loss of efficiency. For EW WZ events, mistiming of ECAL trigger objects in forward and high momentum jets can lead to such L1 pre-firing. The jet leading to pre-firing is incorrectly associated with the previous bunch crossing, producing an event missing the leptonic signature of WZ events. The triggered event is likely to be rejected by the HLT, while the correctly-timed event is rejected by the restriction against successive triggered events.

A correction for this efficiency loss with respect to the expectation in the MC simulation is determined using an unbiased sample of events in which the event pre-firing is prevented by additional trigger conditions. In particular, a sample of events is selected where the trigger conditions restricting the readout of events separated by less than three bunch crossings prevented the mistimed electromagnetic signal from firing the trigger, but did not prevent the correctly timed event from triggering. This sample of events is formed by selecting jet candidates in the reconstructed data, defined using the anti- k_T algorithm and required to have $p_T > 30 \text{ GeV}$, which are matched to their trigger primitive objects. The jet pre-firing efficiency is defined as the ratio of events with jets matched to an L1 trigger object associated with the correct bunch crossing to the all events with selected jets. For jets with $|\eta| < 2.25$, the pre-firing rate is less than 2% over the entire p_T spectrum. In contrast, it grows as a function of jet p_T and η up to values close to 100% for jets with $2.75 < |\eta| < 3.0$ and $p_T > 400 \text{ GeV}$.

The impact of this inefficiency on this analysis is estimated by applying these efficiency corrections, estimated in bins of jet p_T and η , to the selected signal and background events. Due to the signature of forward and high momentum jets, the EW WZ process is particularly affected. This is demonstrated in Fig. 6.2, which shows the estimated efficiency loss in jet η , m_{jj} , and $\Delta\eta(j_1, j_2)$ for simulated EW WZ

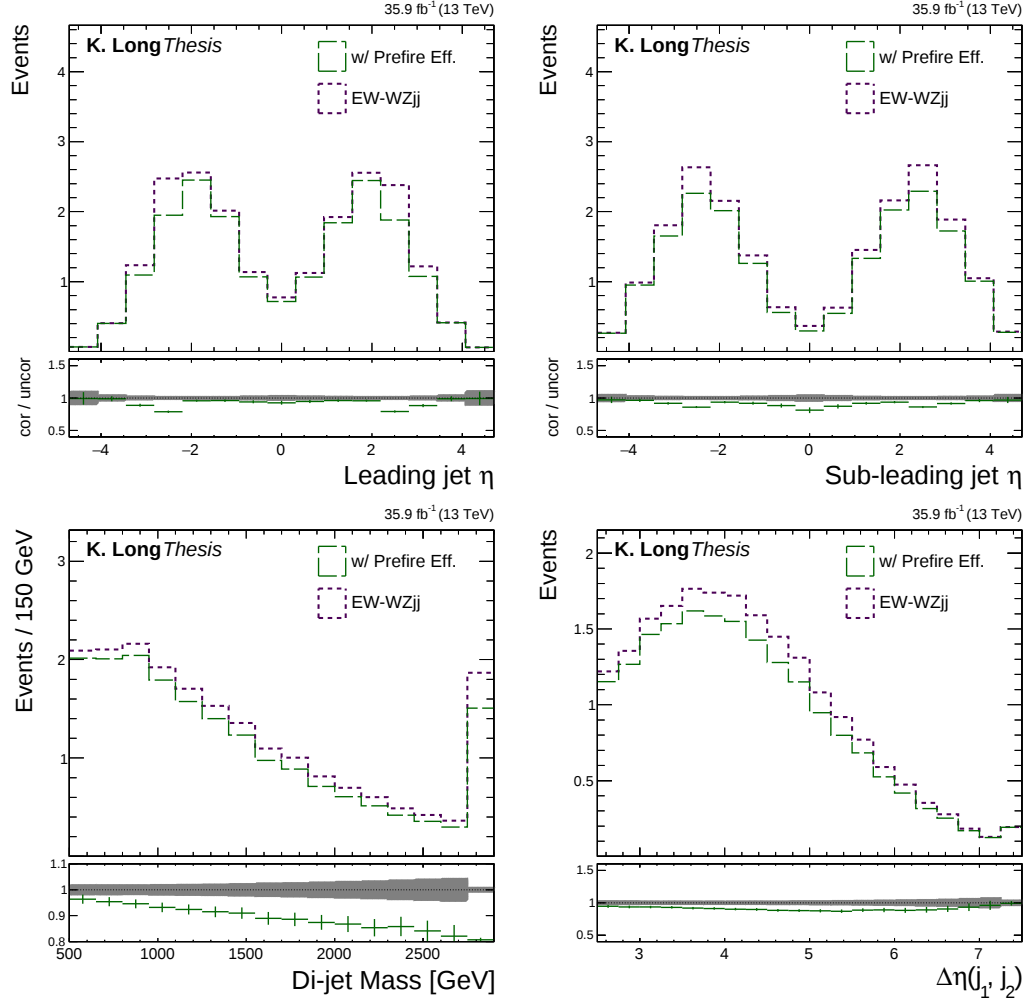


Figure 6.2: Estimated efficiency loss for selected EW WZ simulated events due to mistimed events in the L1 trigger. The solid lines show the uncorrected distributions, while the dotted lines show the efficiency-corrected predictions.

events selected for the analysis. This loss of efficiency is found to be about 1% for m_{jj} of 200 GeV and it increases to about 15% for $m_{jj} > 2$ TeV. The loss of sensitivity due to this inefficiency in the EW WZ search is estimated to be about 10%.

6.3 Event selection and fiducial cross section regions

A selected event is required to have three lepton candidates $\ell\ell'\ell'$, where $\ell, \ell' = e, \mu$. All leptons must pass the identification and isolation requirements described in Chapter 5. The electrons and muons can be directly produced from a W or Z boson decay or from a W or Z boson with an intermediate τ lepton decay. The $\ell'\ell'$ pair consists of two leptons with opposite charge and the same flavor, as expected for a Z boson candidate. One of the leptons from the Z boson candidate is required to have $p_{\text{T}}^{\ell'_1} > 25 \text{ GeV}$ and the other $p_{\text{T}}^{\ell'_2} > 15 \text{ GeV}$. For events with three same-flavor leptons, two oppositely charged, same-flavor combinations are possible. The pair with invariant mass closest to $m_{\text{Z}} = 91.2 \text{ GeV}$, the nominal Z boson mass from Ref. [49], is selected as the Z boson candidate. The remaining lepton is associated with the W boson and must have $p_{\text{T}}^{\ell} > 20 \text{ GeV}$. Events containing additional isolated leptons with $p_{\text{T}}^{\ell} > 10 \text{ GeV}$ are rejected. Because of the neutrino in the final state, the events are required to have $p_{\text{T}}^{\text{miss}} > 30 \text{ GeV}$. To reduce contributions from $t\bar{t}$ events, the leptons constituting the Z boson candidate are required to have an invariant mass satisfying $|m_{\ell'\ell'} - m_{\text{Z}}| < 15 \text{ GeV}$ and events with a b tagged jet with $p_{\text{T}}^{\text{b}} > 30 \text{ GeV}$ and $|\eta^{\text{b}}| < 2.4$ are vetoed.

The invariant mass of any dilepton pair $m_{\ell\ell}$ must be greater than 4 GeV. Such a requirement is necessary in theoretical calculations to avoid divergences from collinear emission of same-flavor opposite-sign dilepton pairs, and 4 GeV is chosen to avoid low mass resonances and double semileptonic decays of B and D hadrons. The selection is extended to all dilepton pairs to reduce contributions from backgrounds with soft leptons while having a negligible effect on signal efficiency. The trilepton invariant mass, $m_{3\ell}$, is required to be more than 100 GeV to exclude a region where production

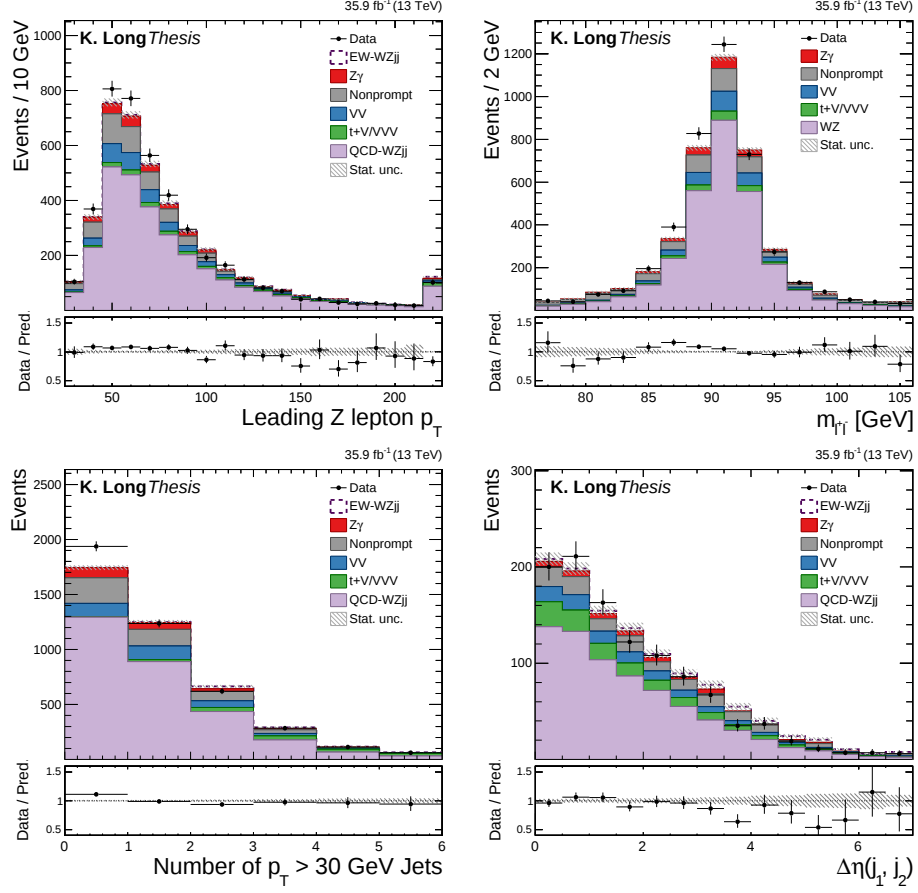


Figure 6.3: Distributions for events satisfying the WZ selection described in the text. The p_T of the highest momentum lepton associated with the Z boson candidate (top left), the invariant mass of the Z candidate (top right), and the event rate by number of anti- k_T jets with $p_T > 30$ GeV (bottom left) are shown with no requirement on the jet activity in the event. The rapidity Separation of the highest p_T jets for events with at least two $p_T > 30$ GeV events are also shown (bottom right).

of Z bosons with final-state photon radiation is expected to contribute. Such events enter the signal region when the radiated photon pair converts to electrons and one electron is lost.

Distributions of events satisfying these criteria, designed to select $WZ \rightarrow 3\ell\nu$ events, are shown in Fig. 6.3. The WZ production rate by number of anti- k_T jets, and the rapidity separation of the two highest p_T jets for events with at least two anti- k_T jets with are also shown.

With no condition on the jet activity in the selection, the EW WZ contribution is negligible. To enhance the EW WZ contribution, the event must have at least two jets with $p_T > 30 \text{ GeV}$ and $|\eta| < 4.7$. The jet with the highest p_T is called the leading jet and the jet with the second-highest p_T the subleading jet. To exploit the unique signature of the VBS process, these two jets are required to have $m_{jj} > 500 \text{ GeV}$ and η separation $|\Delta\eta(j_1, j_2)| \equiv |\Delta\eta_{jj}| > 2.5$. This selection is used for searches for charged Higgs bosons and therefore called the “Higgs boson selection.” For the WZjj and EW WZ measurement, the leading and subleading jet must satisfy $p_T > 50 \text{ GeV}$, and the variable $\eta_{3\ell}^* = \eta^{3\ell} - \frac{1}{2}(\eta^{j_1} + \eta^{j_2})$ of the three-lepton system is additionally required to be between -2.5 and 2.5 . This selection is referred to as the “EW signal selection.” A summary of these selections is shown in Table 6.1.

To reduce the dependence on theoretical predictions, measurements are reported in two fiducial regions, defined in Table 6.1. The “tight fiducial region” is defined to be as close as possible to the measurement phase space, whereas the “loose fiducial region” is designed to be easily reproducible in theoretical calculations or in MC simulations, following the procedure of Ref. [169]. The fiducial predictions are defined through selections on particle-level simulated events using the RIVET [194] framework, which provides a toolkit for analyzing simulated events in a model-independent way. Electrons and muons are required to be prompt (i.e., not from hadron decays), and those produced in the decay of a τ lepton are not considered in the definition of the fiducial phase space. The momenta of prompt photons located within a cone of radius $\Delta R = 0.1$ are added to the lepton momentum to correct for final-state photon radiation, referred to as “dressing.” The three highest p_T leptons are selected and associated with the W and Z bosons with the same procedure used in the data selection. The fiducial cross section in the QCD WZ sideband region is defined following the tight fiducial region of Table 6.1, with $m_{jj} > 100 \text{ GeV}$ and

Table 6.1: Summary of event selections and fiducial region definitions for the analysis. The selections labeled “EW signal” and “Higgs boson” are applied to data and reconstructed simulated events. The EW signal selection is used for all measurements except for the charged Higgs boson search that uses the selection indicated in the column labeled “Higgs boson.” The WZjj cross section is reported in the fiducial regions defined by the selections specified in the last two columns applied to particle-level simulated events. The variables n_j and n_b refer to the number of anti- k_T jets and the number of anti- k_T b-tagged jets, respectively. Other variables are defined in the text.

	EW signal	Higgs boson	Tight fiducial	Loose fiducial
$p_T^{\ell'1}$ [GeV]	>25	>25	>25	>20
$p_T^{\ell'2}$ [GeV]	>15	>15	>15	>20
p_T^ℓ [GeV]	>20	>20	>20	>20
$ \eta^\mu $	<2.4	<2.4	<2.5	<2.5
$ \eta^e $	<2.5	<2.5	<2.5	<2.5
$ m_{\ell\ell} - m_Z $ [GeV]	<15	<15	<15	<15
$m_{3\ell}$ [GeV]	>100	>100	>100	>100
$m_{\ell\ell}$ [GeV]	>4	>4	>4	>4
p_T^{miss} [GeV]	>30	>30	—	—
$ \eta^j $	< 4.7	<4.7	<4.7	<4.7
p_T^j [GeV]	>50	>30	>50	>30
$ \Delta R(j, \ell) $	>0.4	>0.4	>0.4	>0.4
n_j	≥ 2	≥ 2	≥ 2	≥ 2
p_T^b [GeV]	>30	>30	—	—
$ \eta^b $	<2.4	<2.4	—	—
n_b	=0	=0	—	—
m_{jj}	>500	>500	>500	>500
$ \Delta\eta_{jj} $	>2.5	>2.5	>2.5	>2.5
$ \eta^{3\ell} - \frac{1}{2}(\eta^{j1} + \eta^{j2}) $	<2.5	—	<2.5	—

$m_{jj} < 500$ GeV or $|\Delta\eta_{jj}| < 2.5$ or $|\eta_{3\ell}^*| > 2.5$. Theoretical predictions are evaluated using MADGRAPH5_aMC@NLO at LO interfaced to PYTHIA with the samples described in Section 6.4.1.

6.4 Background contributions

Background contributions in this analysis are divided into two categories: background processes with prompt isolated leptons, e.g., ZZ, tZq, $t\bar{t}Z$; and background processes

with nonprompt leptons from hadrons decaying to leptons inside jets or jets misidentified as isolated leptons, primarily $t\bar{t}$ and Z +jets. The background processes with prompt leptons are estimated from MC simulation, while backgrounds with nonprompt leptons from hadronic activity are estimated from control samples in data. The nonprompt component of the $Z\gamma$ process, in which the photon experiences conversion into leptons in the tracker, is evaluated using MC simulation.

6.4.1 Signal simulation and estimation of prompt backgrounds

Several Monte Carlo (MC) event generators are used to simulate the signal and background processes.

The EW-induced production of WZ boson pairs and two final-state quarks, where the W and Z bosons decay leptonically, is simulated at leading order (LO) in perturbative QCD using MADGRAPH5_aMC@NLO v2.4.2 [141]. The sample includes all contributions to the four-lepton final state at $\mathcal{O}(\alpha^6)$ and with a resonant W boson propagator, including triboson processes, where the WZ boson pair is accompanied by a third vector boson that decays into jets, as well as diagrams involving the quartic coupling vertex. The resonant W boson is decayed using MADSPIN [166]. Contributions with an initial-state b quark are excluded from the sample as they are considered part of the tZq process. The predictions from this sample are cross-checked with LO predictions from the event generators VBFNLO 3.0 [173] and SHERPA v2.2.4 [140, 195], and with fixed-order calculations from MoCANLO+RECOLA [142, 169]. Agreement is obtained when using equivalent configurations of input parameters, including couplings, particle masses and widths, and the choice of renormalization (μ_R) and factorization scales (μ_F), and differences arising from such configurations are considered

when assessing the uncertainty.

Several MC simulations of the QCD WZ process are considered. The simulations are inclusive in the number of jets associated with the leptonically decaying W and Z bosons, and therefore include the WZjj contribution studied in this analysis. The primary MC sample is simulated at LO with MADGRAPH5_aMC@NLO v2.4.2, with contributions to WZ production with up to three outgoing partons included in the matrix element calculation. The different jet multiplicities are merged using the MLM scheme [159]. A next-to-LO (NLO) sample from MADGRAPH5_aMC@NLO v2.3.3 with zero or one outgoing partons at Born level, merged using the FxFx scheme [157], and an inclusive NLO sample from POWHEG 2.0 [154, 155, 175, 176] are also considered. The LO MC with MLM merging, referred to as the MLM-merged sample, is taken as the central prediction due its inclusion of WZ plus three-parton contributions at tree level, which are relevant to WZjj production. The other samples, which are used to access the modeling uncertainty in the QCD WZ process, are referred to as the FxFx-merged and the POWHEG samples, respectively. Each sample is scaled to the inclusive NLO cross section from POWHEG 2.0.

The contribution from QCD WZ production is estimated with MC simulation. It is considered signal for the WZjj cross section measurement, but is the dominant background for the EW WZ measurement and in searches for new physics. For the EW WZ measurement and new physics searches, the normalization of the QCD WZ process is constrained by data in the QCD WZ sideband region. The cross section predicted by the MLM-merged sample in the QCD WZ sideband region is $18.6^{+2.9}_{-2.3}(\text{scale}) \pm 1.0(\text{PDF}) \text{ fb}$, where the scale and PDF uncertainties are calculated using the procedure described in Section 6.6. In this region the normalization correction, which is derived from a fit to the data, is consistent with unity. The EW WZ process, considered signal for the WZjj and EW WZ measurements but background

to new physics searches, is also estimated using MC simulation.

In addition to the EW WZ and QCD WZ process, which at tree level are $\mathcal{O}(\alpha^4)$ and $\mathcal{O}(\alpha^2\alpha_S^2)$ respectively, a smaller contribution at $\mathcal{O}(\alpha^3\alpha_S)$ contributes to the WZjj state. We refer to this contribution as the interference term. It amounts to less than 5% of the signal in the EW signal selection and is taken as an uncertainty on the signal contribution. It is evaluated using samples of particle-level events generated with MADGRAPH5_aMC@NLO v2.6.0. Samples are generated with the dynamic μ_R and μ_F set to the maximum of the parton transverse momenta per event, and with fixed scales $\mu_R = \mu_F = m_W$, where m_W is the world average value of the W boson mass, taken from Ref. [49].

The associated production of a Z boson and a single top quark, referred to as tZq production, is simulated at NLO in the four-flavor scheme using MADGRAPH5_aMC@NLO v2.3.3. The sample is scaled using a cross section computed at NLO with MADGRAPH5_aMC@NLO in the five-flavor scheme, following the procedure of Ref. [196].

The production of Z boson pairs via $q\bar{q}$ annihilation is generated at NLO in perturbative QCD with POWHEG 2.0 while the $gg \rightarrow ZZ$ process is simulated at LO with MCFM 7.0 [89]. The ZZ samples are scaled to the cross section calculated at next-to-NLO for $q\bar{q} \rightarrow ZZ$ [197] (K factor 1.1) and at NLO for $gg \rightarrow ZZ$ [198] (K factor 1.7). The $Z\gamma$, $t\bar{t}V$ ($t\bar{t}W$, $t\bar{t}Z$), tZq, and triboson events VVV (WWZ, WZZ, ZZZ) are generated at NLO with MADGRAPH5_aMC@NLO v2.3.3, with the vector bosons generated on-shell and decayed via MADSPIN.

The simulation of the aQGC processes is performed at LO using MADGRAPH5_aMC@NLO v2.4.2 and employs matrix element reweighting to obtain a finely spaced grid of parameters for each of the anomalous couplings operators probed by the analysis. The configuration of input parameters is equivalent to that used for the EW WZ sample described previously. The production of charged Higgs bosons in the Georgi-

Machacek model is simulated using MADGRAPH5_aMC@NLO v2.3.3.

The PYTHIA v8.212 [199, 200] package is used for parton showering, hadronization, and the underlying event simulation, with parameters set by the CUETP8M1 tune [165] for all simulated samples. For the EW WZ process, comparisons are made at particle-level with the parton shower and hadronization of SHERPA and with HERWIG v7.1 [148, 201]. For all MC simulations used in this analysis, the NNPDF3.0 [136] set of parton distribution functions (PDFs) is used, with PDFs calculated to the same order in perturbative QCD as the hard scattering process.

The detector response is simulated using a detailed description of the CMS detector implemented in the GEANT4 package [127, 128]. The simulated events are reconstructed using the same algorithms as used for the data. The simulated samples include additional interactions in the same and neighboring bunch crossings, referred to as pileup. Simulated events are weighted so that the pileup distribution reproduces that observed in the data, which has an average of about 23 interactions per bunch crossing.

6.4.2 Estimation of nonprompt backgrounds

The contribution from background processes with nonprompt leptons is evaluated with control regions of events in data. Events satisfying the full analysis selection, with the exception that one, two, or three leptons pass relaxed identification requirements but fail the more stringent requirements applied to signal events, are selected to form loose lepton control regions. The loose lepton control regions are mutually independent and additionally independent from the signal selection. The small contribution to the loose lepton control regions from events with three prompt leptons is estimated with MC simulation and subtracted from the event samples.

A lepton candidate is labeled as “passed” (P) if it passes the “tight” identification

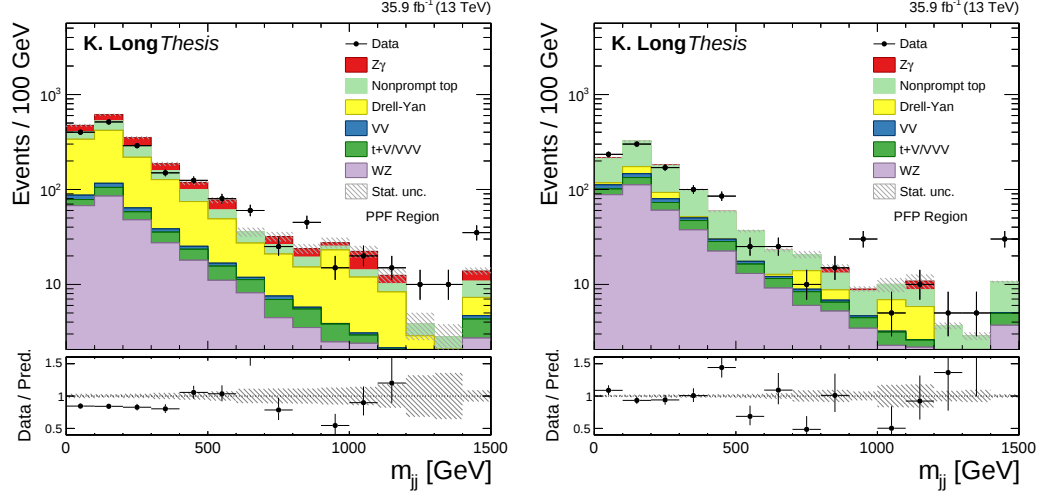


Figure 6.4: The PPF (left) and PFP (right) control regions used in the nonprompt estimation for events satisfying the WZjj selection. The nonprompt MC simulations are shown for reference, and demonstrate that the PPF region is dominated by Drell–Yan events while the PFP region is dominated by $t\bar{t}$ events.

requirements required for signal events, and “failed” (F) if it passes the relaxed lepton identification criteria, but fails the tight identification. The signal region is defined by events with three passing leptons, referred to as PPP events. The loose lepton control regions are defined by the identification criteria of the three selected lepton candidates. The PPF region, where the first two labels refer to the two leptons which satisfy the tight identification criteria and are associated with the Z boson candidate, and the failing lepton is associated with the W-candidate, is dominated by Z+jet events. The m_{jj} for the observed data compared to the prediction of the MC simulation in the WZjj selection, with $p_T^j > 30$ GeV, is shown in two illustrative control regions in Fig. 6.4. The nonprompt estimation does not depend on the nonprompt MC simulations, including Drell–Yan and $t\bar{t}$ production, which are shown for reference. All possible combinations of P and F objects are listed in Table 6.2, along with the dominant production mechanisms of such events.

The expected contribution from these processes in the signal region is estimated

Table 6.2: Possible combinations of lepton identification for three-lepton events that define the control regions used to estimate nonprompt backgrounds.

Dominant contribution	Combination		
	Z-candidate leptons	W-candidate lepton	
QCD multi-jet	F	F	F
W+jets	F	F	P
	P	F	F
	F	P	F
Z+jets	P	P	F
$t\bar{t}$, other	P	F	P
	F	P	P

using “loose-to-tight” efficiency factors, which are applied as extrapolation factors to the lepton candidates failing the analysis requirements in the control region events. The efficiency factors, referred to as “fake rates,” are evaluated using samples of events dominated by jet activity. A sample of $Z+\ell_{\text{candidate}}$ events is defined by selecting events with exactly three leptons satisfying the loose identification criteria, using the same trigger requirements as for signal events. Two opposite sign, same flavor leptons satisfying the tight identification and consistent with the decay of a Z boson, specifically $|m_{\ell^+\ell^-} - m_Z| < 10 \text{ GeV}$, where m_Z denotes the nominal Z boson mass, are required. To reduce contamination from WZ events, the events must additionally have $p_T^{\text{miss}} < 25 \text{ GeV}$ and $m_T(\ell_3, p_T^{\text{miss}}) < 25 \text{ GeV}$. The lepton not associated with the Z boson candidate is taken as the $\ell_{\text{candidate}}$ object.

The fake rates are evaluated using ratios of events where the $\ell_{\text{candidate}}$ object satisfies the full identification requirements to events where all identification criteria are not satisfied, defined as $\varepsilon_{\text{fake}} = \frac{N_{\text{tight}}}{N_{\text{loose}}}$. They are parameterized as a function of the $\ell_{\text{candidate}}$ p_T and η . The expected contribution from true three-prompt-lepton events is estimated from MC simulation and subtracted from the event samples. The fake rates measured in data, corrected using the prompt MC simulation, and the predicted fake rates from the nonprompt MC simulation are in good agreement, as shown in

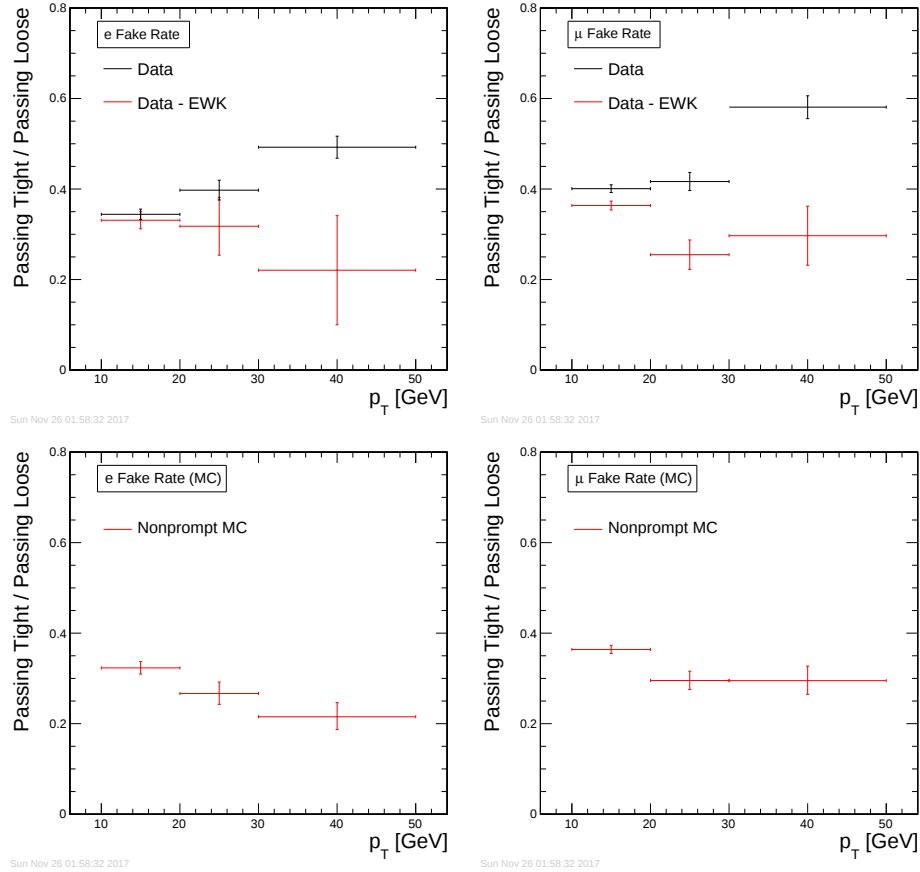


Figure 6.5: Ratio of events in which the $\ell_{\text{candidate}}$ objects passes the tight lepton identification criteria to candidates passing only the loose lepton identification criteria in bins of the $\ell_{\text{candidate}}$ p_T for the selection of $Z+\ell_{\text{candidate}}$ events described in the text. The events in which the $\ell_{\text{candidate}}$ object is a electron (muon) are shown in the left (right) plot. The top plots show events selected in data, with corrections for contamination of true prompt events from MC simulation. The bottom plots show predictions from the $Z+\text{jet}$ and $t\bar{t}$ MC samples.

Fig. 6.5, The MC simulation prediction is not used in the analysis, but serves as a consistency-check of the nonprompt estimation technique. The two-dimensional values of the fake rates derived from data which are used for the analysis are shown in Fig. 6.6.

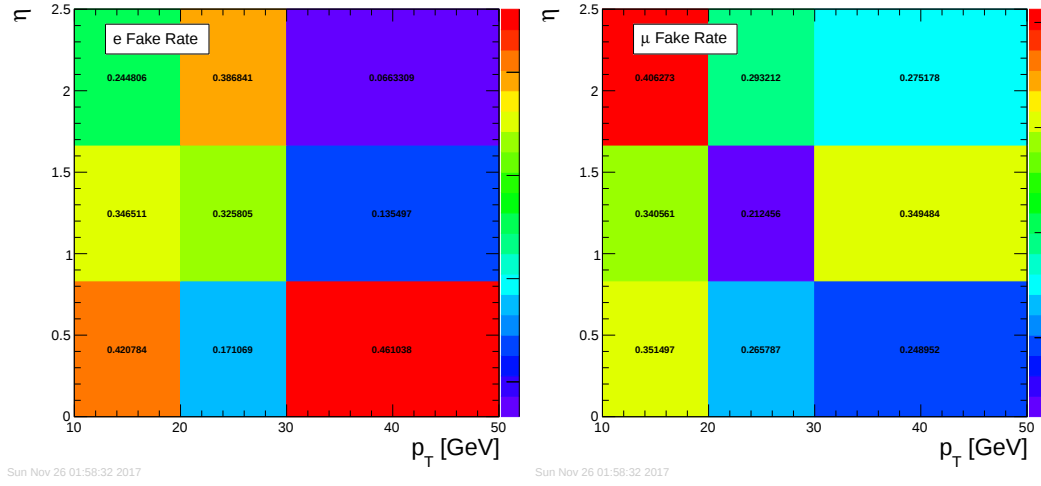


Figure 6.6: Ratio of events in data, with corrections for contamination of true prompt events from MC simulation, in which the $\ell_{\text{candidate}}$ objects passes the tight lepton identification criteria to candidates passing only the loose lepton identification criteria in bins of p_T and η for the selection of $Z+\ell_{\text{candidate}}$ events described in the text. The events in which the $\ell_{\text{candidate}}$ object is a electron (muon) are shown in the left and right plots respectively.

A cross-check of the technique is performed by repeating the procedure with efficiency factors derived from a sample of events dominated by dijet production. Events passing single lepton triggers are selected, requiring at least one jet and one $\ell_{\text{candidate}}$ object passing the loose identification. Events with additional loose leptons are not selected. The loose lepton must be separated from the jet with highest p_T by $\Delta R(\ell_{\text{candidate}}, j) > 1$. The contamination of WZ is suppressed by requiring $E_T^{\text{miss}} < 25$ GeV and $m_T < 25$ GeV. The loose-to-tight efficiency factors obtained in the two regions agree to within 30% for the full p_T and η range.

Because control regions with n_F failing leptons can contribute to the region with

Control region	Includes regions						
FPP	FPP			FFF	FPP	FFP	
PFP			PFP	FFF		FFP	PFF
PPF		PPF		FFF	FPP		PFF
FPP				FFF	FPP		
FFP				FFF		FFP	
PFF				FFF			PFF
FFF				FFF			

Table 6.3: Possible nonprompt control regions and regions of overlap.

$n_F - 1$ failing leptons, the total background contribution in the signal region is not a simple sum of the loose lepton control region event yields multiplied by the fake rate of each failing leptons. Table 6.3 shows the possible contributions of loose lepton control regions with loose leptons into tight lepton regions. For example, the estimated contribution of events to the signal region $N_{B,FPP+PFP}$ from the FPP and PFP loose lepton control regions, with N_{FPP} and N_{PFP} data events is given by

$$\begin{aligned}
N_{B,FPP+PFP} = & N_{FPP} \times f_{l1} - \underbrace{N_{FFF} \times f_{l1} \times f_{l2} \times f_{l3}}_{\text{interference of FPP and FFF}} \\
& - \underbrace{f_{l1} \times f_{l3} \times (N_{FPP} - N_{FFF} \times f_{l2})}_{\text{interference of FPP and FPP}} - \underbrace{f_{l1} \times f_{l2} \times (N_{FPP} - N_{FFF} \times f_{l3})}_{\text{interference of FPP and FPP}} \\
& + N_{PFP} \times f_{l2} - \underbrace{N_{FFF} \times f_{l1} \times f_{l2} \times f_{l3}}_{\text{interference of PFP and FFF}} - \underbrace{f_{l1} \times f_{l2} \times (N_{PFP} - N_{FFF} \times f_{l3})}_{\text{interference of PFP and FFP}} \\
& - \underbrace{f_{l2} \times f_{l3} \times (N_{PFP} - N_{FFF} \times f_{l1})}_{\text{interference of PFP and PFF}} = \\
& N_{FPP} \times f_{l1} + N_{PFP} \times f_{l2} - 2N_{FFF} \times f_{l1} \times f_{l2} \times f_{l3} - 2N_{FPP} \times f_{l1} \times f_{l2} \\
& - N_{FPP} \times f_{l1} \times f_{l3} - N_{PFP} \times f_{l2} \times f_{l3}
\end{aligned}$$

where $f_{l1} = \frac{\varepsilon_{\text{fake}}}{1 - \varepsilon_{\text{fake}}}$, computed from the misidentification ratio $\varepsilon_{\text{fake}}$ in the particular p_T and η bin.

Combining the yields from all seven loose lepton control regions in this way yields

the following formula for the total number of estimated background events from all regions

$$N_B = N_{FPP} \times FR_{l1} + N_{PFP} \times FR_{l2} + N_{PPF} \times FR_{l3} + N_{FFF} \times FR_{l1} \times FR_{l2} \times FR_{l3} \\ - N_{FPP} \times FR_{l1} \times FR_{l3} - N_{PFP} \times FR_{l1} \times FR_{l2} - N_{PPF} \times FR_{l2} \times FR_{l3}.$$

The characteristics of the control region events are used to obtain per-channel differential predictions in the signal region. The expected contribution of nonprompt events in the signal region is shown in Fig. 6.9, where the signal and background normalizations are fixed to their expected values and the expected sum of signal and background events is treated as the observed data.

This method is validated in nonoverlapping data samples enriched in Drell–Yan and $t\bar{t}$ contributions. The Drell–Yan region is defined by inverting the selection requirement in p_T^{miss} , and the $t\bar{t}$ region is defined by requiring at least one b-tagged jet and rejecting events with $|m_{\ell'\ell'} - m_Z| < 5 \text{ GeV}$ while keeping all other requirements for the signal region. The background prediction and observed data in these regions for events satisfying the WZjj selection are shown in Fig. 6.7, which demonstrates agreement between the two techniques. The predictions for large values of m_{jj} and $\Delta\eta_{jj}$ are limited by the event count of data events used for the nonprompt estimation.

The event count of the loose lepton control samples of nonprompt events satisfying the EW signal selection severely limits differential predictions in this region. This is illustrated in Fig. 6.8, which shows the differential prediction from the nonprompt background in the EW signal region, which is particularly limited when dividing distributions by decay channel.

The combined shape of the nonprompt background for all channels is therefore used for each channel in the EW WZ measurement and in the extraction of constraints

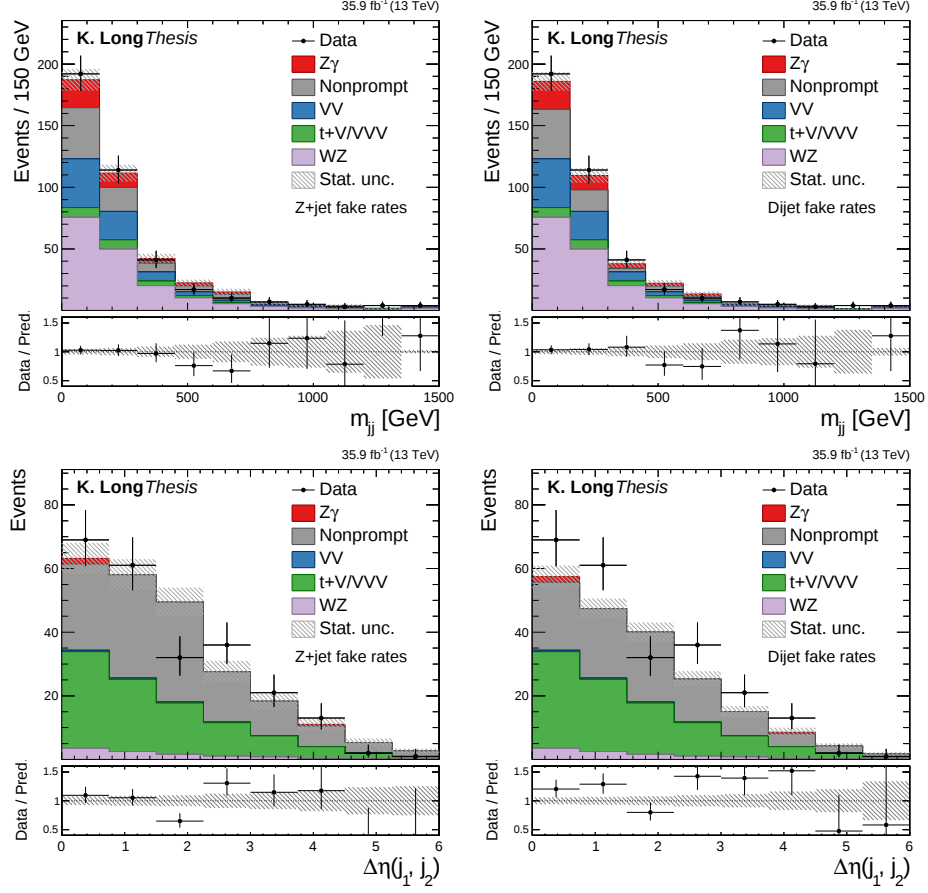


Figure 6.7: m_{jj} for events in the Drell–Yan control region (top) and $\Delta\eta_{jj}$ for events in the $t\bar{t}$ control region (bottom). The events are required to have two $p_T > 30$ GeV jets but no additional kinematic selection is applied to the dijet system.

on anomalous quartic gauge couplings. The normalization of the distribution per channel is taken from the ratio of the nonprompt yield in a single channel to the total nonprompt event yield measured in WZjj events with no requirements on the dijet system. These ratios are found to be consistent within the statistical uncertainty with ratios measured when relaxing the jet p_T requirement in WZjj events, in WZ events inclusive in the number of jets, and in events satisfying the EW signal and control region selections, as demonstrated in Table 6.4.

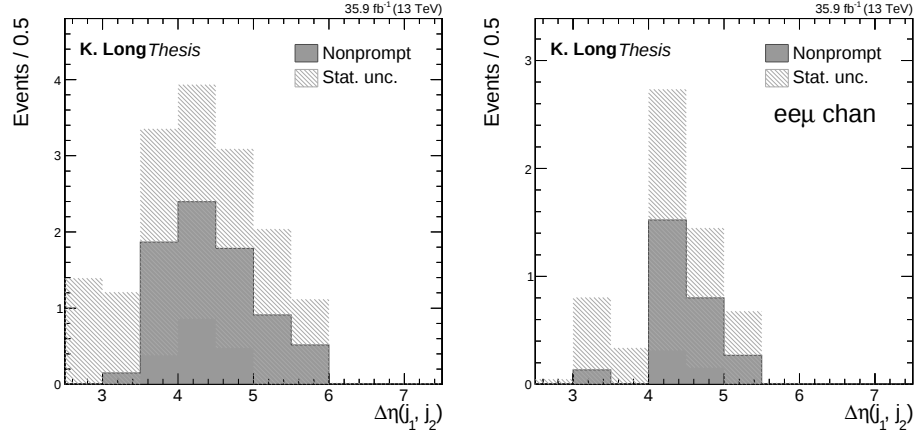


Figure 6.8: Distribution of $\Delta\eta_{jj}$ for estimated nonprompt events in the EW signal region. The left plot shows the combined contribution for all channels, while the right plot shows the contribution from only the $ee\mu$ channel, which is severely limited by the event count in the loose lepton control samples.

Table 6.4: Relative fraction of nonprompt events by channel for different selections of WZ and WZjj events.

Selection	eee	ee μ	$\mu\mu e$	$\mu\mu\mu$
WZ inclusive	0.06 ± 0.04	0.18 ± 0.03	0.21 ± 0.04	0.55 ± 0.04
WZjj	0.08 ± 0.05	0.16 ± 0.05	0.33 ± 0.07	0.44 ± 0.06
Charged Higgs boson	0.14 ± 0.14	0.22 ± 0.14	0.30 ± 0.21	0.34 ± 0.15
EW signal	$< 0.1 \pm 0.1$	0.23 ± 0.19	0.57 ± 0.34	0.21 ± 0.16

6.5 Statistical proceduress

The distribution of data events in a quantum process is fundamentally stochastic in nature. For scattering experiments with a low expected number of events, the probability $f(n, \nu)$ for an experiment to observe n events when measuring a quantity with an expected event count ν is given by the Poisson distribution,

$$f(n, \nu) = e^{-\nu} \frac{\nu^n}{n!}. \quad (6.1)$$

The expected events arise from both signal and background processes, that is, $\nu = \nu_b + \nu_s$, where ν_b and ν_s represent the expected event yields from background

and signal processes respectively. As this analysis seeks to extract information from the data about the signal processes—which may differ from the expectation—the two contributions are not treated on equal footing. The variability of the signal expectation is quantified by introducing a rate modifier μ for the expected signal event yield, which is referred to as the “signal strength” for the signal process.

Given this relationship, an observation of n events combined with an estimation of ν_b can be used to estimate the value for μ . The estimated value of the signal process production cross section σ_s follows trivially, given the relationship between the integrated luminosity $\mathcal{L}$ of the scattering process, the cross section, and the expected event yield, $\nu_s = \mathcal{L}\sigma_s$.

The probability distribution in terms of the observed and expected event yields can therefore be expressed as

$$f(n, \nu) = e^{-(\nu_b + \mu\sigma_s\mathcal{L})} \frac{(\nu_b + \mu\sigma_s\mathcal{L})^n}{n!}. \quad (6.2)$$

Given the stochastic nature of the observation n , the true value of ν_b and ν_s from which the data are drawn cannot be inferred precisely but can be estimated by leveraging knowledge of the underlying distribution of the event yields. Uncertainties in the underlying model present an additional complication in quantifying the probability of an observation in terms of an input model, as they reduce the accuracy in which the underlying distribution is known. In this analysis, the knowledge of expected number of background-associated events is limited by several limitations in the model, including the energy scale of jet measurements, the theoretical uncertainty in the WZ processes and other contributions discussed in Section 6.6. Estimating the expected event distributions and their uncertainties forms the central challenge of this analysis, as discussed in the previous sections.

The expected probability distribution is therefore a function of additional nuisance parameters, which themselves have a distribution of possible values. That is,

$$f(n, \nu, \vec{\theta}) = e^{-\nu(\vec{\theta})} \frac{\nu(\vec{\theta})^n}{n!}, \quad (6.3)$$

where the dependence of the expected event yield on parameters of the model $\vec{\theta}$ is made explicit in this expression.

The probability distribution of the expected event yields is particularly important when a measurement seeks to quantify the consistency of the observed data with a spectrum of possible signal hypotheses. Of particular importance is the “background-only hypotheses,” in which $\mu = 0$. For a given observed n , one may seek to quantify whether the observation is unambiguously consistent with the $\mu \neq 0$ or if the data can fully be described by the background expectation.

The statistical procedure to quantify this consistency in most particle physics experiments is the maximum likelihood estimation, or maximum likelihood fit. The technique seeks to estimate the underlying parameters of the model from which the data set is drawn, using knowledge of the expected distribution of the model parameters. The technique is built around the likelihood function, which is built from the of the expected probability distributions for the outcome of an experiment. The likelihood is a function of the parameters of the model which are not known with certainty. Given an observed data set, the best estimate for the true set of parameters of the model from which the data maximizes this function [202].

Results presented in this analysis use a binned likelihood function, which is a product of the probability distribution functions per bin. That is,

$$\mathcal{L}(n, \nu, \vec{\theta}) = \prod_{i=0}^{n_{bins}} e^{-\nu_i(\vec{\theta})} \frac{\nu_i(\vec{\theta})^{n_i}}{n_i!} \times e^{-\theta_i^2/2}. \quad (6.4)$$

In this expression, the set of parameters $\vec{\theta}$, referred to as nuisance parameters, accounts for the uncertainty in the expected values of the associated model parameter. Each $\vec{\theta}$ parameter has a probability density function around its central value which is derived for each parameter and discussed in Section 6.6. As shown in Equation 6.4, as $\vec{\theta}$ parameters diverge from their input values $\theta_i = 0$, corresponding to no variation from the expected central value of a parameter, the likelihood function is reduced. The parameter may impact the predictions in all bins equally, referred to as a normalization uncertainty, or it may affect the distribution in a bin-dependent way, referred to as a shape uncertainty. For the extraction of results in this analysis, log-normal probability density functions are assumed for the nuisance parameters affecting the event yields of the various background contributions, whereas systematic uncertainties that affect the shape of the distributions are represented by nuisance parameters whose variation results in a continuous perturbation of the spectrum [203] and which are assumed to have a Gaussian probability density function.

The interpretations of results in terms of the SM predictions and BSM extensions presented in this thesis are carried out via binned maximum likelihood fits of the predicted yields and distributions to the observed data. The form of the binned likelihood used in each analysis interpretation is presented briefly in the following discussion. Fits to the likelihood are used to quantifying the statistical significance of an observed excess over the background-only hypothesis, to extract results in terms of cross section measurements, and establish one- or multi-dimensional confidence intervals for BSM model parameters. The likelihood function is maximized, or rather the function $-\ln \mathcal{L}$, using the MINUIT software package in the ROOFIT framework. Distributions and expected event yields shown using the best-fit normalizations and uncertainties from this procedure are referred to as post-fit distributions, whereas predictions representing the input values of parameters to fit are referred to as pre-fit

results.

6.5.1 Fiducial WZjj cross section measurement

For the extraction of the fiducial WZjj cross sections, the likelihood function is built from the predicted yields in the four leptonic channels for the EW signal selection, shown in Equation 6.5. A single signal strength for WZjj production μ_{WZjj} is treated as an unconstrained parameter in the fit. Explicitly,

$$\mathcal{L}(\mu, \vec{\theta} | n_c) = \prod_c^{\text{channels}} e^{-(\nu_{b,c}(\vec{\theta}) + \mu_{\text{WZjj}} \sigma_{th} \times \mathcal{B}_{\text{WZ} \rightarrow c} \mathcal{L})} \frac{(\nu_{b,c}(\vec{\theta}) + \mu_{\text{WZjj}} \sigma_{th} \times \mathcal{B}_{\text{WZ} \rightarrow c} \mathcal{L})^{n_c}}{n_c!} e^{-\theta^2/2}. \quad (6.5)$$

The likelihood is a function of the number of expected background events ν_b , which is a function of the nuisance parameters $\vec{\theta}$, which may be channel dependent. Here σ_{th} is the theoretical prediction for the WZjj cross section, $\mathcal{B}_{\text{WZ} \rightarrow c}$ is the branching ratio for WZ to each of the four leptonic decay channels, and $\mathcal{L}$ is the integrated luminosity. The expectation parameters used at input to the likelihood are shown in Fig. 6.9, where the expected sum of signal and background predictions is treated as the observed data in the fit.

The observed cross section is obtained from the value of μ_{WZjj} which maximizes the likelihood function given

$$\sigma_{\text{WZjj,obs}} = \mu_{EW} \times \sigma_{th}. \quad (6.6)$$

Note that this expression does not imply that $\sigma_{\text{WZjj,obs}}$ depends directly on the predicted cross section, as μ_{WZjj} will adjust linearly to fit the observed data, as seen in Equation 6.5. Thus any input value for σ_{th} will lead to the same $\sigma_{\text{WZjj,obs}}$ given by Equation 6.6, provided the efficiency for reconstructing the events is unaffected.

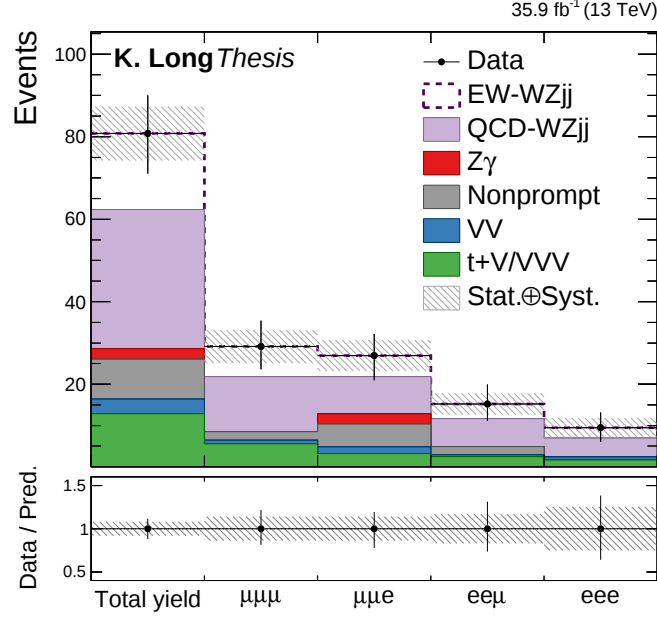


Figure 6.9: Expected composition of signal and background events for the EW signal selection. The expected sum of signal and background contributions are treated as the observed data in a maximum likelihood fit. The uncertainty in the total expected event yield from WZ and non-WZ contributions is shown as a hashed band.

In terms of the efficiency for reconstruction ϵ , which is given by the ratio of expected events in the fiducial volume to the number of reconstructed events, $\epsilon = n_{th}/n_{reco}$.

In this expression, the number of expected events n_{th} excludes events outside the fiducial volume of the detector, or outside the kinematic region where events can reasonably be reconstructed, e.g., leptons at very low p_T . For this reason, the cross section obtained from this expression is referred to as a fiducial cross section. To report a cross section in a region that difference significantly from the experimental fiducial volume, the extrapolation factor from the region of measurement must be calculated. This fully theoretical quantity, known as the acceptance, introduces a theoretical uncertainty to the measurement. It is defined as $\mathcal{A} = \sigma_{report,th}/\sigma_{meas,th}$ where $\sigma_{meas,th}$ and $\sigma_{report,th}$ are the predicted cross sections in the region of measurement and reported region respectively. The fiducial cross section regions considered

in this analysis are described in Section 6.3.

6.5.2 Search for EW WZ boson production

Separating the EW- and QCD-induced components of WZjj events requires exploiting the different kinematic signatures of the two processes. The relative fraction of the EW WZ process with respect to the QCD WZ process and other backgrounds grows with increasing values of the m_{jj} and $|\Delta\eta_{jj}|$ the leading jets, as demonstrated in Fig. 6.10, which shows the expected distributions of m_{jj} and $\Delta\eta_{jj}$ for events satisfying the EW signal selection. This motivates the use of a 2D distribution built from these variables for the extraction of the EW WZjj signal strength via a maximum likelihood fit. This 2D distribution, shown as a one-dimensional histogram in Fig. 6.11, along with the yield in the control region defined in Section 6.3, are combined in a binned likelihood involving the expected and observed numbers of events in each bin. The distribution including the yield in the QCD WZ control region, shown as the first bin, is shown in Fig. 6.12.

The likelihood is a combination of individual likelihoods for the four decay channels, following the definition of Equation 6.4. The measurement seeks to compare the best-fit value of μ_{EW} from the maximum likelihood, to the background-only fit in which $\mu_{EW} = 0$. The significance of μ_{EW} is a measure of the compatibility of the measurement with the best-fit value of $\mu_{EW} > 0$ obtained from the maximum likelihood fit, referred to as the signal hypothesis, to the background-only hypothesis.

The compatibility of the measurement with the background-only hypotheses is quantified via the test statistic, t_μ , which is a function of the best-fit signal strength

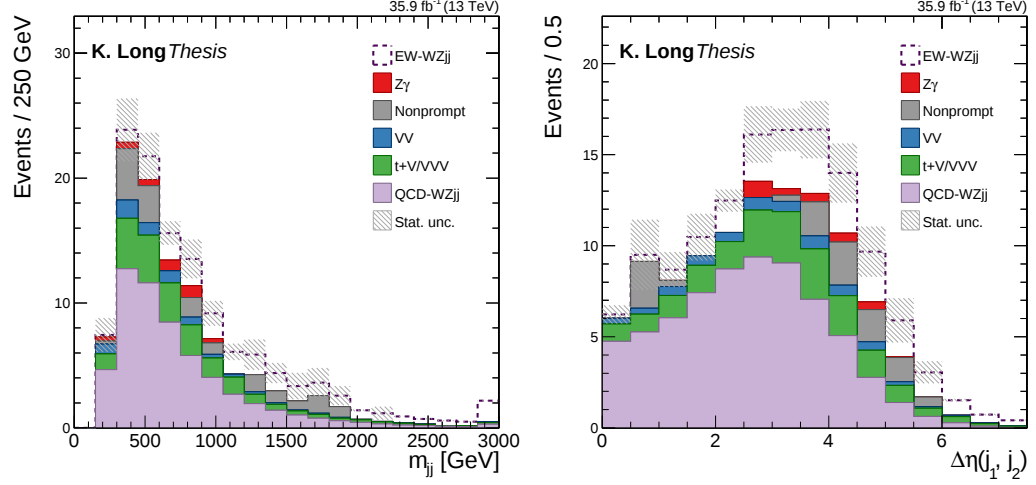


Figure 6.10: The m_{jj} (left) and $|\Delta\eta_{jj}|$ of the two leading jets (right) for events satisfying the EW signal selection. The last bin contains all events with $m_{jj} > 2500$ GeV (left) and $|\Delta\eta_{jj}| > 7.5$ (right). The dashed line shows the expected EW WZ contribution stacked on top of the backgrounds, which are shown as filled histograms. The hatched bands represent the total and relative statistical uncertainties on the predicted yields. The bottom panel shows the ratio of the number of events measured in data to the total number of expected events. The predicted yields are shown with their pre-fit normalizations.

and the model with $\mu = 0$. It is defined as

$$t_\mu = \left\{ \begin{array}{ll} \frac{\mathcal{L}(\mu, \hat{\theta}(\mu))}{\mathcal{L}(\hat{\mu}, \hat{\theta})} & \mu \geq \hat{\mu} \\ \frac{\mathcal{L}(\mu, \hat{\theta}(\mu))}{\mathcal{L}(0, \hat{\theta})} & \mu < 0 \end{array} \right\}, \quad (6.7)$$

where $\hat{\mu}$ and $\hat{\theta}$ are the values of μ and $\hat{\theta}$ which maximize the likelihood, and $\hat{\theta}(\nu)$ are the values of $\hat{\theta}$ which maximize the likelihood for the given μ . The significance of a measured value $\hat{\mu} \neq 0$ is quantified via the test statistic $q_0 = t_\mu(\mu = 0)$. Note that the consistency with the background-only hypothesis is evaluated with respect to the signal plus background hypothesis, and only a positive signal strength is considered a signal-like deviation from the background.

The test statistic has a probability distribution $f(q_\mu|\mu)$ for a given μ . This distribution can be determined by sampling the underlying model and nuisance parameters

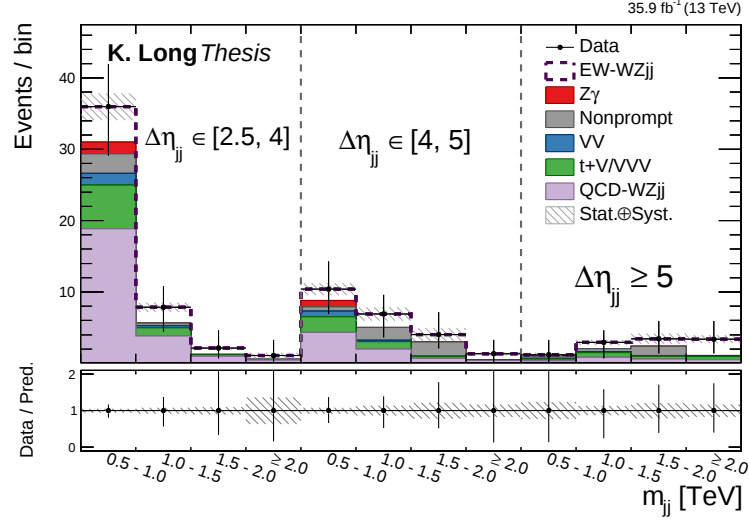


Figure 6.11: The one-dimensional representation of the 2D distribution of m_{jj} and $|\Delta\eta_{jj}|$, used for the EW signal extraction. The x axis shows the dijet mass distribution in the indicated bins, split into three bins of $\Delta\eta_{jj}$: $\Delta\eta_{jj} \in [2.5, 4]$, $[4, 5]$, ≥ 5 . The dashed line represents the EW WZ contribution stacked on top of the backgrounds, which are shown as filled histograms. The hatched bands represent the total and relative systematic uncertainties on the predicted yields. The bottom panel shows the ratio of the number of events measured in data to the total number of expected events. The predicted yields are shown with their best-fit normalizations.

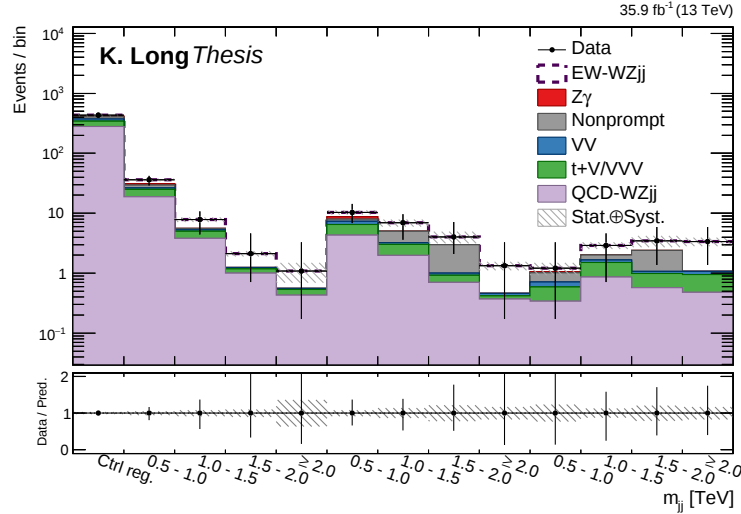


Figure 6.12: Visualization of the likelihood function used for the EW WZ search. The first bin represents the yield in the low- m_{jj} control region, with other bins showing the one-dimensional representation of 2D distribution of m_{jj} and $|\Delta\eta_{jj}|$ shown in 6.11.

to form an ensemble of pseudo-experiments drawn from a value of μ . The observed value $\hat{\mu}$ obtained from a fit to the data corresponds to a single realisation of the possible observations given the underlying model. As sampling of the distribution is extremely computationally extensive, it is common for particle physics results to be derived with asymptotic formulae to approximate the distribution of $f(q_\mu|\mu)$ following Ref. [204]. These formulae are applicable provided the observed number of events per bin in the likelihood is not small. All results in this thesis use the asymptotic approach.

The probability of this result given the expected parameters of the model in the background-only hypothesis is quantified in terms of the p -value,

$$p_\mu = \int_{t_{\mu,obs}}^{\infty} f(t_\mu|\mu) dt. \quad (6.8)$$

The value of p_0 (or p -value) is customarily transposed such that the area from $x = \pm\sigma$ to $x = \pm\infty$ for a Gaussian function with unit area and standard deviation σ would give the same value. This value, in terms of the standard deviation, is referred to as the significance of the result. It is customary in the particle physics community to refer to a p -value of 2.87×10^{-7} , corresponding to a significance of 5 standard deviations, or 5σ , as sufficient for rejection of the background-only hypothesis in favor of the signal plus background model.

The expected significance of EW WZ process obtained from a fit to the likelihood shown in Fig. 6.12, based on the expected distribution of background and signal events with $\mu_{EW} = 1$, is 2.5 standard deviations.

6.5.3 Limits on charged Higgs boson production

For the EW WZ search, the background-only assumption corresponds to the SM with no EW WZ process. An excess with respect to the background-only hypothesis is

therefore expected, and is quantified using the procedures of the previous section. For searches for BSM physics, one does not anticipate the need to quantify the significance of a signal-like fluctuation. One seeks instead to place constraints on the parameters of the unobserved BSM model based on the consistency of the search with the null hypothesis.

Upper limits on the signal strength μ are derived using the CL_s criterion of Refs. [204–206]. The procedure follows the approach of the previous section, with a slight modification of the test statistic, defined as

$$t_\mu = \left\{ \begin{array}{ll} \frac{\mathcal{L}(\mu, \hat{\theta}(\mu))}{\mathcal{L}(\hat{\mu}, \hat{\theta})} & \mu \geq \hat{\mu} \\ \frac{\mathcal{L}(\mu, \hat{\theta}(\mu))}{\mathcal{L}(0, \hat{\theta})} & \mu < \hat{\mu} \end{array} \right\}. \quad (6.9)$$

The distribution of expected results and the p -value of the observation are calculated from this test statistic for a range of possible values of μ for the signal hypothesis using Equation 6.8. The CL_s criterion is defined as

$$\text{CL}_s = \frac{p_\mu}{p_0}. \quad (6.10)$$

The signal strength μ that satisfies the condition $\text{CL}_s = 1 - \alpha$ is said to be excluded at confidence level (CL) α . For example, the value $\alpha = 0.95$ corresponds to an exclusion limit of 95% CL.

Because presence of charged Higgs bosons would give rise to a resonance-like structure in the WZ transverse mass, defined as

$$m_{\text{T}}(\text{WZ}) = \sqrt{(E_{\text{T}}(\text{W}) + E_{\text{T}}(\text{Z}))^2 - (\vec{p}_{\text{T}}(\text{W}) + \vec{p}_{\text{T}}(\text{Z}))^2}, \quad (6.11)$$

with $E_{\text{T}} = \sqrt{m^2 + p_{\text{T}}^2}$, where the W candidate is constructed from the $\vec{p}_{\text{T}}^{\text{miss}}$ and the lepton associated with the W boson, the likelihood used to test the consistency of the data with the presence of a charged Higgs boson is built from this distribution,

shown in Fig. 6.13, and the event yield in the control region described in Section 6.4.1. A combined fit of the predicted signal and background yields to the data in the Higgs boson selection is performed to derive model-independent expected and observed upper limits on $\sigma_{\text{VBF}}(\text{H}^\pm) \mathcal{B}(\text{H}^\pm \rightarrow \text{WZ})$ at 95% CL using the CL_s criterion.

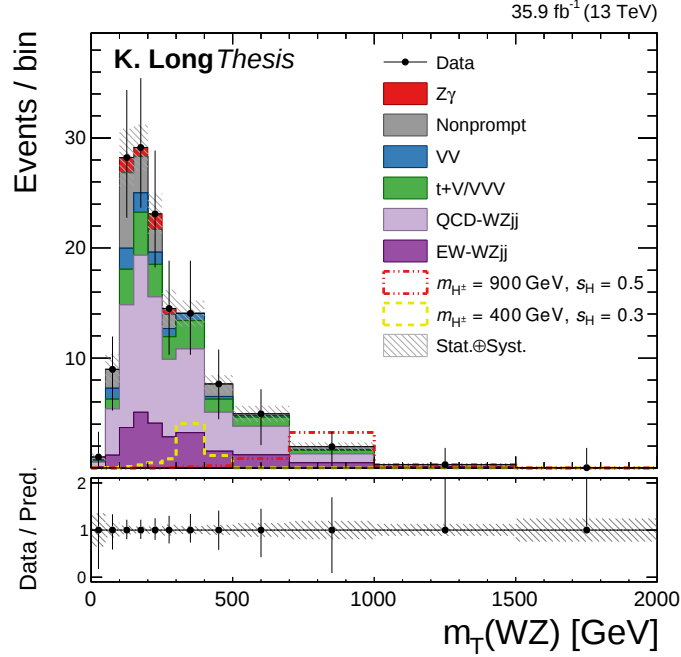


Figure 6.13: The $m_T(\text{WZ})$ for events satisfying the Higgs boson selection, used to place constraints on the production of charged Higgs bosons. The last bin contains all events with $m_T(\text{WZ}) > 2000$ GeV. The dashed lines show predictions from the Georgi-Machek model with $m(\text{H}^\pm) = 500$ (900) GeV and $s_H = 0.3$ (0.5). The bottom panel shows the ratio of the number of events measured in data to the total number of expected events. The hatched bands represent the total and relative systematic uncertainties on the predicted background yields. The predicted yields are shown with their best-fit normalizations from the background-only fit.

6.5.4 Limits on anomalous quartic gauge couplings

Events satisfying the EW signal selection are used to constrain aQGCs in the effective field theory approach [207]. Results are obtained following the formulation of Ref. [79], which proposes nine independent dimension-eight operators which assume

the presence of a light Higgs boson and respect the $SU(2) \times U(1)$ symmetry of the EW gauge sector. All operators are charge conjugation and parity-conserving. The $WZjj$ channel is most sensitive to the $T0$, $T1$, and $T2$ operators, which are constructed purely from the $SU(2)$ gauge fields, the $S0$ and $S1$ operators, which involve interactions with the Higgs field, and the $M0$ and $M1$ operators, which involve a mixture of gauge and Higgs field interactions.

The presence of nonzero aQGCs would enhance the production of events with high WZ mass. The binned likelihood used for the extraction of results is therefore built from $m_T(WZ)$ as for the limits on charged Higgs bosons. The expected distribution is shown in Fig. 6.14, with several illustrative anomalous coupling parameters also shown.

The MC simulations of non-zero aQGCs include the SM EW WZ process, with an increase in the yield at high $m_T(WZ)$ arising from parameters different from their SM values. Because the increase of the expected yield over the SM prediction exhibits a quadratic dependence on the operator coefficient, a parabolic function is fitted to the predicted yields per bin to obtain a smooth interpolation between the discrete operator coefficients considered in the MC simulation. The one-dimensional 95% CL limits are extracted using the CL_s criterion as for the charged Higgs boson search. The SM prediction, including the EW WZ process, is treated as the null hypothesis.

One-dimensional limits are placed with all parameters except for the coefficient being probed set to zero. In addition, a scan of the likelihood is also performed with multiple independent μ values corresponding to different dimension-eight operators.

Constraints are also placed on aQGC parameters using a two-dimensional scan, where two parameters are probed in the fit with all others set to zero. In this case, the likelihood is a function of two signal strength parameters, and a two-dimensional exclusion region is defined by the region in μ_1 and μ_2 space in which the likelihood

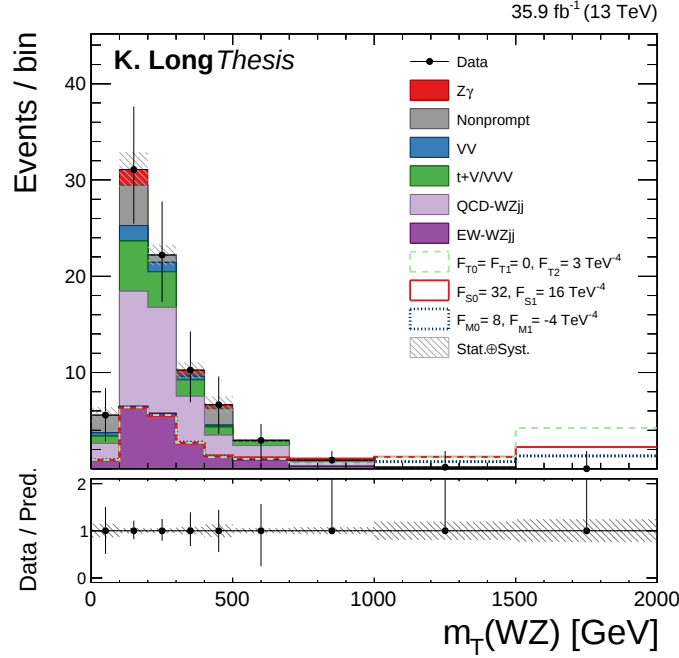


Figure 6.14: The $m_T(\text{WZ})$ for events satisfying the EW signal selection, used to place constraints on the anomalous coupling parameters. The dashed lines show predictions for several aQGC parameters values, which modify the EW WZ process. The last bin contains all events with $m_T(\text{WZ}) > 2000 \text{ GeV}$. The hatched bands represent the total and relative systematic uncertainties on the predicted yields. The bottom panel shows the ratio of the number of events measured in data to the total number of expected events. The predicted yields are shown with their best-fit normalizations from the background-only fit.

differs from its maximum value by the desired confidence interval.

6.6 Systematic uncertainties

The dominant uncertainties in both the cross section measurement and new physics searches are those associated with the jet energy scale (JES) and resolution (JER). The JES and JER uncertainties are applied in simulated events by smearing and scaling the relevant observables and propagating the effects to the event selection and the kinematic variables used in the analysis. The uncertainty in the event yield in the EW signal selection due to the JES and JER is found to be 9% for QCD WZ

and 5% for EW WZ processes. The uncertainty additionally depends on the p_T and η of the selected jets. For the QCD WZ process, it varies in the range of 5–25% with increasing values of m_{jj} and $|\Delta\eta_{jj}|$.

Uncertainties in signal and background processes estimated with MC simulation are evaluated from the theoretical uncertainties of the predictions. Event weights in the simulated samples are used to evaluate variations of the central prediction. Scale uncertainties are estimated by varying μ_R and μ_F by a factor of two from their nominal values, with the condition that $1/2 \leq \mu_R/\mu_F \leq 2$. The maximal and minimal variations are obtained per bin to form a shape-dependent variation band. The PDF uncertainties are obtained from the MC replica sets of the NNPDF3.0 PDF. The PDF and scale uncertainties are uncorrelated for different signal and background process and 100% correlated across bins for the distributions used to extract results.

The uncertainty in modeling the EW WZ and QCD WZ processes has a large impact in the EW WZ measurement. In addition to the uncertainties from scale and PDF choice, comparisons of alternative matrix element and parton shower generators are considered. The uncertainty in the QCD WZ process is derived by comparing the predictions of the MLM-merged simulation and those obtained with the FxFx-merged simulation, after fixing the normalization to the observed data in the QCD WZ sideband region. Differences between the predictions of the MC simulations in the signal region and in the ratio of the QCD WZ sideband to the signal region event yields are considered in the comparisons. The differences in predictions are generally within the scale and PDF uncertainties of the MC simulations, and a 10% normalization uncertainty is assigned to account for the observed discrepancies. The results obtained using the POWHEG simulation, which predicts a slightly softer m_{jj} spectrum, are also largely contained within the theoretical uncertainties considered. However, because WZjj events from this simulation arise from soft radiation from

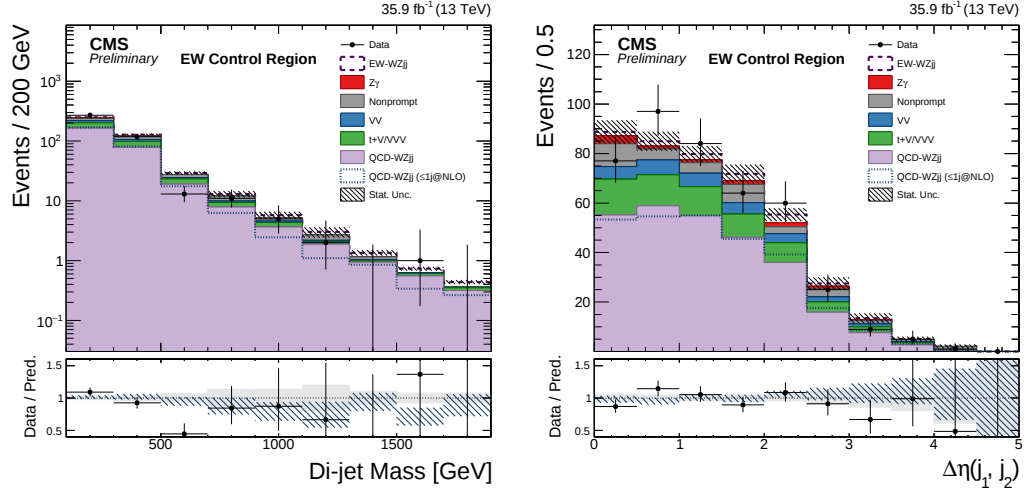


Figure 6.15: The m_{jj} (left) and $\Delta\eta_{jj}$ (right) for events in the non-VBS control region. The last bin contains all events with $m_{jj} > 1900$ GeV (left) and $|\Delta\eta_{jj}| > 5.0$ (right). An alternative prediction for the QCD WZ contribution, using MADGRAPH5_aMC@NLO with FxFx merging as described in the text of SMP-18-001, is shown in dashed blue. This distribution can be compared to the nominal prediction in filled purple. The bottom panel shows the ratio of the number of events measured in data to the total number of expected events for the nominal QCD WZ prediction. The shaded band at 1 represents the size of the statistical uncertainties on the predicted yields. The hatched blue band shows the sum of Monte Carlo predictions and statistical uncertainties with the alternative QCD WZ prediction divided by the sum of Monte Carlo predictions with the nominal QCD WZ prediction. Normalizations are shown as values used for the input to the fit.

the parton shower, it is not explicitly considered in the uncertainty evaluation. For the EW WZ process, the MC simulations described in Section 6.4.1 agree within the theoretical uncertainties from the PDF and the choice of μ_R and μ_F for the kinematic variables considered in the analysis, so no additional uncertainty is assigned.

The interference term is evaluated using MC simulation of particle-level events selected from the samples described in Section 6.4.1. It is found to be positive, and amounts to 12% of the EW WZ contribution in the control region and 4% in the signal region for both MC samples considered. The ratio of the interference to the EW WZ decreases with increasing m_{jj} , consistent with the observations of Ref. [169]. These values are used as a symmetric shape uncertainty in the signal cross section

when performing the EW WZ measurement. This uncertainty is subdominant with respect to other theoretical uncertainties and has a negligible contribution to the uncertainty on the observed EW WZ signal strength.

Higher-order EW corrections in VBS processes are known to be negative and at the level of tens of percent, with the correction increasing in magnitude with increasing m_{jj} [145] and m_{VV} . We do not apply corrections to the WZjj MC simulation, but we have verified that the significance of the EW WZ measurement is insensitive to higher-order EW corrections by performing the signal extraction described in Section 6.5.2 with the m_{jj} predicted by the EW WZ MC simulation modified by the corrections from Ref. [172]. As the relative effect of the EW corrections on SM and anomalous WZjj production is unknown, we do not apply corrections to the SM backgrounds or new physics signals for our results. Because corrections to the SM WZjj production that decrease the expected number of events at high m_{WZ} lead to more stringent limits on new physics, this is a conservative approach.

The uncertainties related to the finite number of simulated events, or to the limited number of events in data control regions, affect the signal and background predictions. They are uncorrelated across different samples, and across bins of a single distribution. The limited number of events in the relaxed lepton control samples used for the nonprompt background estimate is the dominant contribution to this uncertainty.

The nonprompt background estimate is also affected by systematic uncertainties from the jet flavor composition of the control regions and loose-to-tight transfer factors. The systematic uncertainty in the nonprompt event yield is 30% for both electrons and muons, uncorrelated between channels. It covers the largest difference observed between the estimated and measured numbers of events in data control samples enriched in $t\bar{t}$ contributions, which have enhanced contributions from heavy-flavor jets, and Drell–Yan contributions, dominated by light-flavor jets, and the differences

between using transfer factors derived in Z+jet and dijet events.

Systematic uncertainties are less than 1% for the trigger efficiency and 1–3% for the lepton identification and isolation requirements, depending on the lepton flavors. Other systematic uncertainties are related to the use of simulated samples: 1% for the effects of pileup and 1–2% for the E_T^{miss} reconstruction, which is estimated by varying the energies of the PF objects within their uncertainties. The uncertainty in the b-quark jet content in WZ events is 2% and accounts for differences in b-tagging efficiencies between data and MC. The uncertainty in the integrated luminosity of the data sample is 2.5% [CMS:2017sdi]. This uncertainty affects both the signal and the simulated portion of the background estimation, but does not affect the background estimation from data.

Table 6.5: The dominant uncertainty contributions in the fiducial WZjj cross section measurement and their expected contributions to the significance of the EW WZ measurement. The impact of each systematic uncertainty in the WZjj cross section measurement is obtained by freezing the set of associated nuisance parameters to their best fit values and comparing the total uncertainty in the signal strength to the result from the nominal fit. The effect on the EW WZ significance, shown in the last column, is defined as the relative increase in the expected significance when freezing the nuisance term to its best fit value.

Source of systematic uncertainty	Relative systematic uncertainty [%]	
	σ_{WZjj}	EW WZ significance
Jet energy scale	+11 / – 8.1	7.0
Jet energy resolution	+1.9 / – 2.1	< 0.1
QCD WZ modeling	N/A	2.2
Other background theory	+2.2 / – 2.2	0.3
Nonprompt normalization	+2.5 / – 2.5	0.3
Nonprompt event count	+6.0 / – 5.8	1.7
Lepton energy scale and eff.	+3.5 / – 2.7	< 0.1
b tagging	+2.0 / – 1.7	< 0.1
Integrated luminosity	+3.6 / – 3.0	< 0.1

Chapter 7

Results

7.1 Fiducial WZjj cross section measurement

The cross section for WZjj production, without separating by production mechanism, is measured with a combined maximum likelihood fit to the observed event yields by leptonic decay channel of the W and Z bosons as described in Section 6.5. A signal strength μ_{WZjj} , which represents the ratio of the measured signal yield to the expected number of signal events, is treated as a free parameter in the fit.

The best fit value for the signal strength is used to obtain a cross section in the tight fiducial region defined in Table 6.1. The measured fiducial WZjj cross section in this region is

$$\sigma_{\text{WZjj}}^{\text{fid}} = 3.18_{-0.52}^{+0.57} (\text{stat}) \, {}_{-0.36}^{+0.43} (\text{syst}) = 3.18_{-0.63}^{+0.71} \text{ fb}. \quad (7.1)$$

This result can be compared to the predicted value of $3.27_{-0.32}^{+0.39}(\text{scale}) \pm 0.15 (\text{PDF}) \text{ fb}$. The EW WZ and QCD WZ contributions are calculated independently from the samples described in Section 6.4.1 and their uncertainties are combined in quadrature. The interference term contribution in this region is less than 1% of the total cross

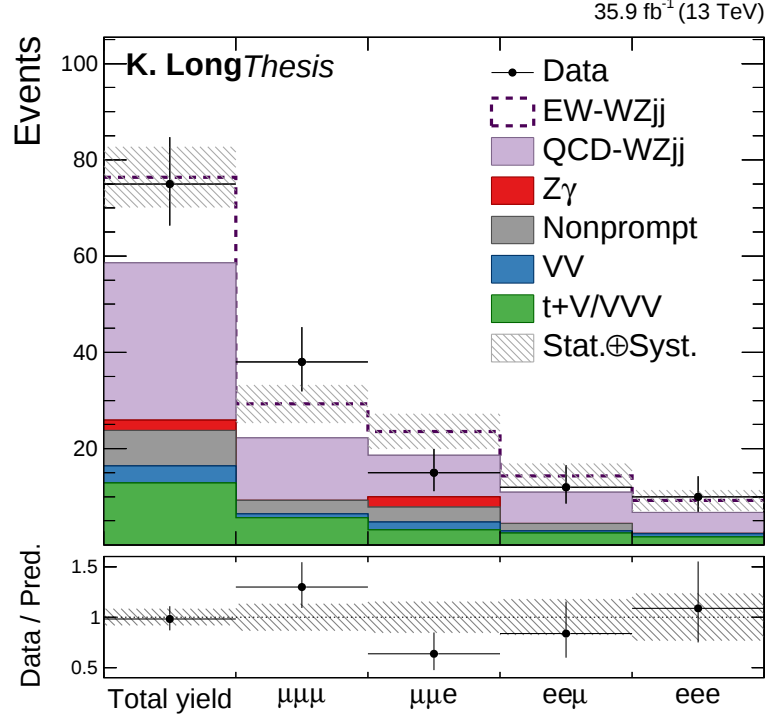


Figure 7.1: Post-fit event yields in the EW signal region.

section.

Results are also obtained in a looser fiducial region, defined in Table 6.1 following Ref. [169], to simplify comparisons with theoretical calculations. The acceptance from the loose to tight fiducial region is $(72.4 \pm 0.8)\%$, computed using MADGRAPH5_aMC@NLO interfaced to PYTHIA. The uncertainty in the acceptance is evaluated independently for the EW WZ and QCD WZ samples from the scale. The uncertainty in the acceptance is evaluated by combining the scale and PDF uncertainties in the EW WZ and QCD WZ predictions in quadrature. The scale uncertainty in the QCD WZ contribution is the dominant component of the uncertainty. The resulting WZjj loose fiducial cross section is

$$\sigma_{\text{WZjj}}^{\text{fid,loose}} = 4.39_{-0.72}^{+0.78} (\text{stat}) \, {}_{-0.50}^{+0.60} (\text{syst}) = 4.39_{-0.87}^{+0.98} \text{ fb}, \quad (7.2)$$

which can be compared to the predicted value of $4.51_{-0.45}^{+0.59}$ (scale) ± 0.18 (PDF) fb. The EW WZ and QCD WZ contributions and their uncertainties are treated independently with the same approach as described for the tight fiducial region.

7.2 Search for EW WZ boson production

The rate of EW WZ production with respect to the SM prediction is quantified by treating the EW WZ process as signal and the QCD WZ process background in a binned likelihood fit to a two-dimensional distribution of m_{jj} and $\Delta\eta(j_1, j_2)$ as described in Section 6.5. The observed data and the background and signal predictions, adjusted to their best-fit values from the maximum likelihood fit, are shown for this distribution in Fig. 7.3. The individual distributions of m_{jj} and $\Delta\eta(j_1, j_2)$ are shown in Fig. 7.2.

The best-fit value for the signal strength μ_{EW} is

$$\mu_{EW} = 0.82_{-0.43}^{+0.51}, \quad (7.3)$$

consistent with the SM expectation at LO of $\mu_{EW, LO} = 1$,

The significance of the signal is quantified by calculating the local p -value for an upward fluctuation of the data relative to the background prediction using a profile likelihood ratio test statistic and asymptotic formulae [204]. The observed (expected) statistical significance for EW WZ production is 2.2 (2.5) standard deviations. The total uncertainty of the signal strength measurement is dominated by the statistical uncertainty of the data. The predicted cross section in the loose fiducial region used in the fit is $1.48_{-0.11}^{+0.12}$ (scale) ± 0.07 (PDF) fb. The post-fit yields for the signal and background corresponding to the best-fit signal strength for EW WZ production are shown in Table 7.1.

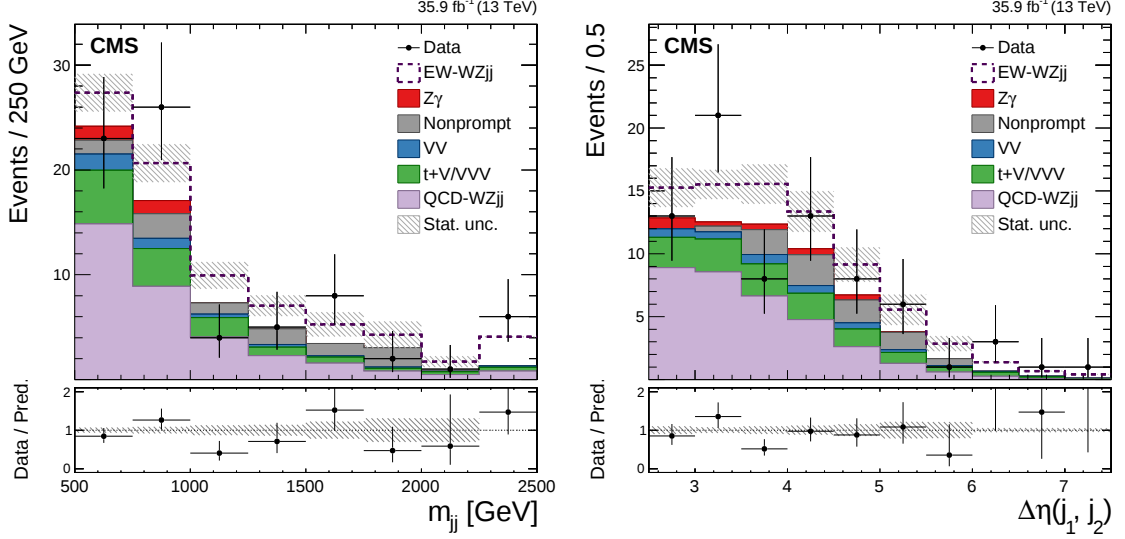


Figure 7.2: The m_{jj} (left) and $|\Delta\eta_{jj}|$ (right) of the two leading jets for events satisfying the EW signal selection. The last bin contains all events with $m_{jj} > 2500$ GeV (left) and $|\Delta\eta_{jj}| > 7.5$ (right). The dashed line shows the expected EW WZ contribution stacked on top of the backgrounds that are shown as filled histograms. The hatched bands represent the total and relative statistical uncertainties on the predicted yields. The bottom panel shows the ratio of the number of events measured in data to the total number of expected events. The predicted yields are shown with their pre-fit normalizations.

Table 7.1: Post-fit event yields after the signal extraction fit to events satisfying the EW signal selection. The EW WZ process is corrected for the observed value of μ_{EW} .

Process	$\mu\mu\mu$	$\mu\mu e$	$ee\mu$	eee	Total yield
QCD WZ	13.5 ± 0.8	9.1 ± 0.5	6.8 ± 0.4	4.6 ± 0.3	34.1 ± 1.1
t+V/VVV	5.6 ± 0.4	3.1 ± 0.2	2.5 ± 0.2	1.7 ± 0.1	12.9 ± 0.5
Nonprompt	5.2 ± 2.0	2.4 ± 0.9	1.5 ± 0.6	0.7 ± 0.3	9.8 ± 2.3
VV	0.8 ± 0.1	1.6 ± 0.2	0.4 ± 0.0	0.7 ± 0.1	3.5 ± 0.2
$Z\gamma$	0.3 ± 0.1	1.2 ± 0.8	<0.1	0.6 ± 0.2	2.2 ± 0.8
Pred. background	25.5 ± 2.1	17.4 ± 1.5	11.2 ± 0.8	8.3 ± 0.6	62.4 ± 2.8
EW WZ signal	6.0 ± 1.2	4.2 ± 0.8	2.9 ± 0.6	2.1 ± 0.4	15.1 ± 1.6
Data	38	15	12	10	75

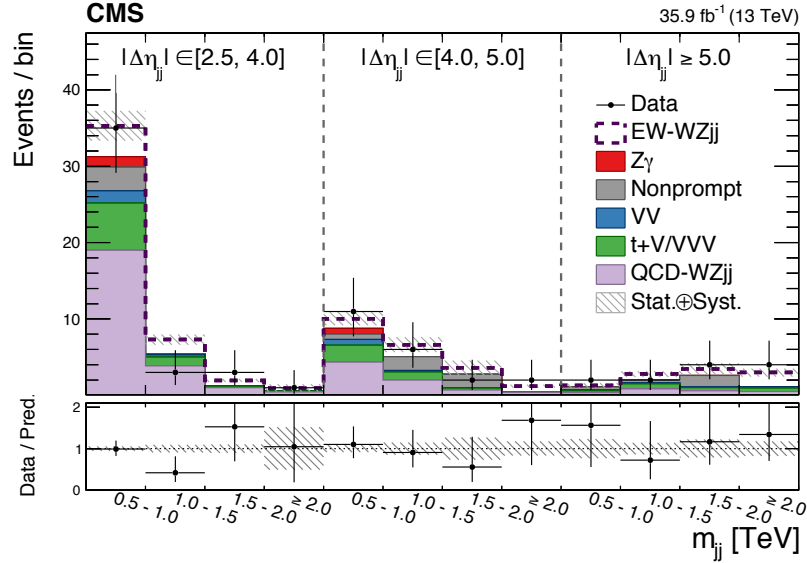


Figure 7.3: The one-dimensional representation of the 2D distribution of m_{jj} and $|\Delta\eta_{jj}|$, used for the EW signal extraction. The x axis shows the m_{jj} distribution in the indicated bins, split into three bins of $\Delta\eta_{jj}$: $\Delta\eta_{jj} \in [2.5, 4]$, $[4, 5]$, ≥ 5 . The dashed line represents the EW WZ contribution stacked on top of the backgrounds that are shown as filled histograms. The hatched bands represent the total and relative systematic uncertainties on the predicted yields. The bottom panel shows the ratio of the number of events measured in data to the total number of expected events. The predicted yields are shown with their best-fit normalizations.

7.2.1 Comparison with the EW WZ measurement by the ATLAS Collaboration

There is a small degree of tension between this result and the result recently submitted for publication by the ATLAS Collaboration [98], which reports observed signal strengths for the EW WZ and QCD WZ processes of $1.77^{+0.49}_{-0.43}$ and 0.56 ± 0.16 , respectively. Because the independent results presented by the ATLAS and CMS Collaborations are each limited by large statistical uncertainties, establishing their mutual compatibility and compatibility with the SM will be aided by future experimental studies. Because the measured signal strength is a ratio of the predicted to observed yields, it is dependent on the MC simulation used to derive the expected

yields. While the cross sections derived from these results are not dependent on the normalization, the fiducial regions used by the two analyses are not directly comparable. Thus, understanding the simulations used in the two independent analyses in a common fiducial region will provide valuable insight into the agreement of the two results. In addition, future measurements with reduced uncertainties will definitively determine if the differences are a result of a statistical fluctuation or if they are a consequence of different analysis techniques or interpretations.

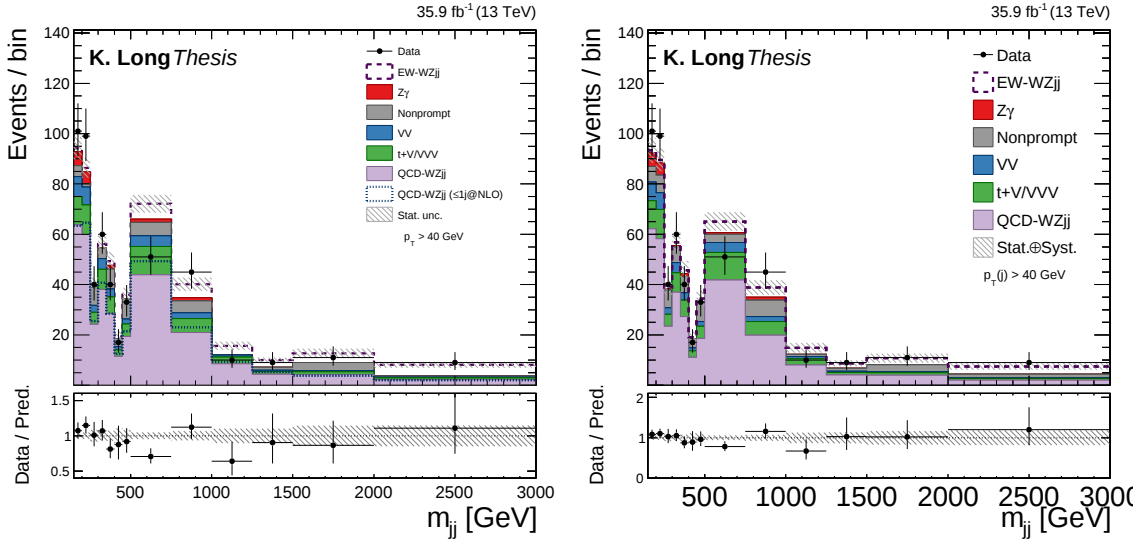


Figure 7.4: The m_{jj} for events selected in a relaxed fiducial region, described in the text, designed to be comparable to the selection used by the ATLAS Collaboration in Ref. [98]. The distribution is shown with the expected normalizations input to the maximum likelihood fit (left) and the with the best fit normalizations by bin obtained from the fit (right). The data are well described by the shape and normalization predicted by the MC simulations of the EW WZ (dashed purple) and QCD WZ (filled light purple) processes discussed in Chapter 4.

The analysis presented in this thesis is designed to be sensitive to the EW WZ process with a minimal and explicit theoretical dependence on well-studied distributions. The approach is emphasized in Fig. 7.4, which shows the m_{jj} for events in a relaxed selection intended to mimic the fiducial region defined by the ATLAS Collaboration. The selection is defined following Table 6.1, with no condition on the $\Delta\eta_{jj}$ or $\eta_{3\ell}^*$ of the

jets and with the jet transverse momentum relaxed to 40 GeV. To validate the results from the nominal EW WZ signal extraction, a maximum likelihood fit is performed to this distribution with the approach described in Chapter 6. While the distribution is less sensitive to the EW WZ process than the two-dimensional likelihood, the central value of the EW WZ signal strength is within 10% of the value presented, well within the uncertainty of the measurement. The normalization correction of the QCD WZ process that maximizes the likelihood is consistent with the presented result and with unity. The distribution is well-described by the MC simulation and background estimation, as illustrated in Fig. 7.4, which shows the distributions input to the maximum likelihood fit and the resulting best-fit distributions. In contrast to the results of Ref. [98], the distribution is expected normalizations for the EW WZ and QCD WZ processes are highly consistent with the observed data. Studies to determine if a multivariate discriminant, following the approach of the ATLAS Collaboration, results in a different categorization of the events as EW WZ and QCD WZ will be a topic of future studies.

7.3 Limits on anomalous quartic gauge couplings

Events satisfying the EW signal selection are used to constrain aQGCs in the effective field theory approach [207]. Results are obtained following the formulation of Ref. [79] via a maximum likelihood fit to the $m_T(\text{WZ})$ as described in Section 6.5.4. The $m_T(\text{WZ})$ for events satisfying the EW signal selection is shown in Fig 7.5. The predictions of several indicative aQGC operators and coefficients are also shown.

The one-dimensional 95% confidence level (CL) limits are extracted using the CL_s criterion [204–206], with all parameters except for the coefficient being probed set to zero. The SM prediction, including the EW WZ process, is treated as the

null hypothesis. No deviation from the SM prediction is observed, and the resulting observed and expected limits are summarized in Table 7.2.

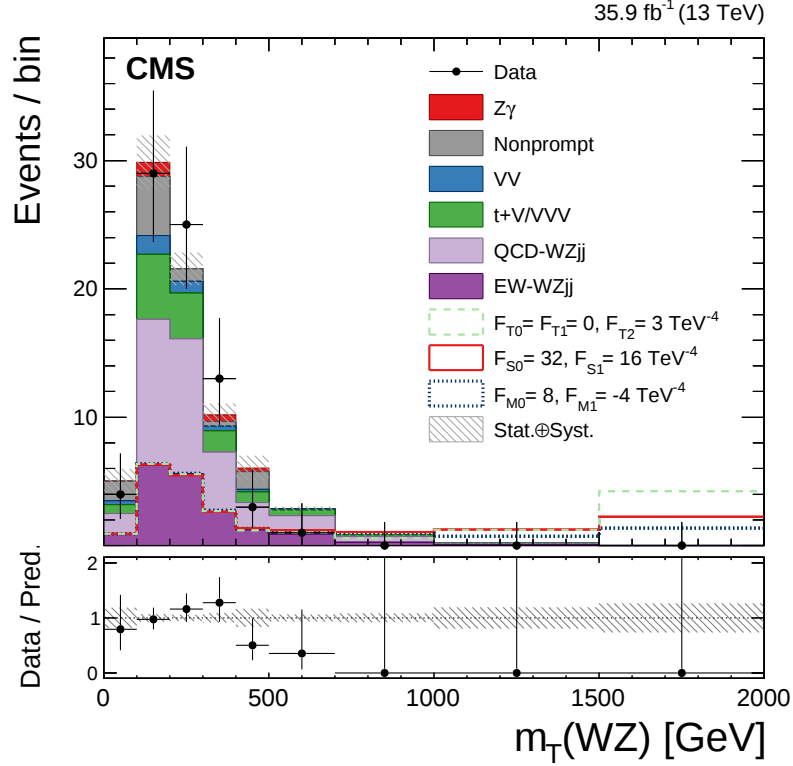


Figure 7.5: The $m_T(\text{WZ})$ for events satisfying the EW signal selection, used to place constraints on the anomalous coupling parameters. The dashed lines show predictions for several aQGC parameters values that modify the EW WZ process. The last bin contains all events with $m_T(\text{WZ}) > 2000$ GeV. The hatched bands represent the total and relative systematic uncertainties on the predicted yields. The bottom panel shows the ratio of the number of events measured in data to the total number of expected events. The predicted yields are shown with their best-fit normalizations from the background-only fit.

Constraints are also placed on aQGC parameters using a two-dimensional scan, where two parameters are probed in the fit with all others set to zero. The resulting 2D 95% CL intervals for these parameters are shown in Fig. 7.6.

Constraints on aQGCs in terms of dimension-eight effective field theory operators have also been obtained by analyses of EW VV events in other production and decay channels. The results presented here are competitive with constraints from analyses

Table 7.2: Observed and expected 95% CL limits for each operator coefficient while all other parameters are set to zero.

Parameters	Expected limit (TeV^{-4})	Observed limit (TeV^{-4})
f_{M0}/Λ^4	$[-11.2, 11.6]$	$[-9.15, 9.15]$
f_{M1}/Λ^4	$[-10.9, 11.6]$	$[-9.15, 9.45]$
f_{S0}/Λ^4	$[-32.5, 34.5]$	$[-26.5, 27.5]$
f_{S1}/Λ^4	$[-50.2, 53.2]$	$[-41.2, 42.8]$
f_{T0}/Λ^4	$[-0.87, 0.89]$	$[-0.75, 0.81]$
f_{T1}/Λ^4	$[-0.56, 0.60]$	$[-0.49, 0.55]$
f_{T2}/Λ^4	$[-1.78, 2.00]$	$[-1.49, 1.85]$

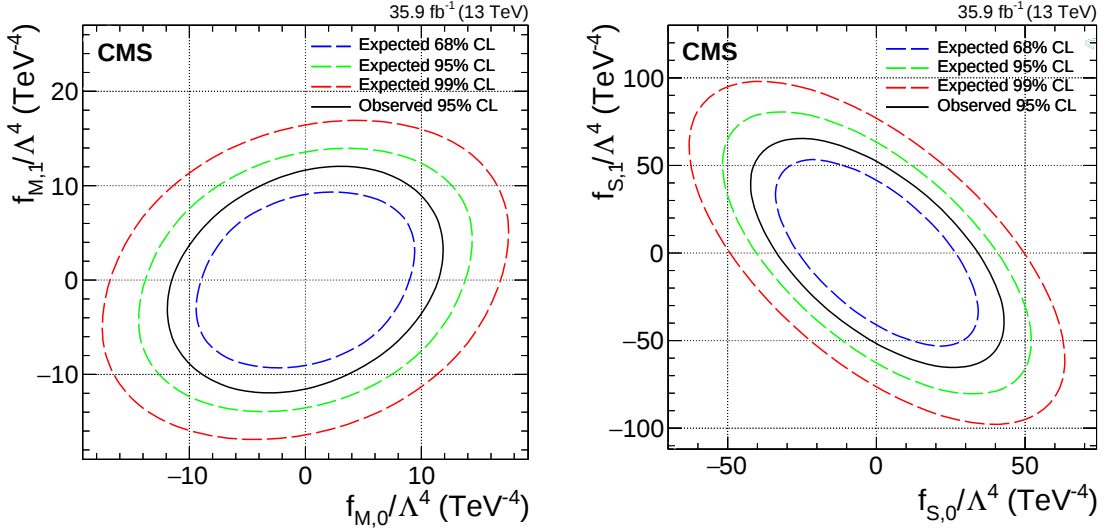


Figure 7.6: Two-dimensional observed 95% CL intervals (solid contour) and expected 68, 95, and 99% CL intervals (dashed contour) on the selected aQGC parameters. The values of coefficients outside of contours are excluded at the corresponding CL.

exploiting leptonic decays, as shown in Fig. 7.8 and 7.7 for the $T0$, $T1$, $T2$, $S0$, and $S1$ operators. The most stringent results using leptonic decays are obtained by the CMS Collaboration using the $W^\pm W^\pm$ [94] channel at 13 TeV. A preliminary study using VV events where one boson decays hadronically and one leptonically has recently been performed by the CMS Collaboration [208], which significantly extends the constraints on operators sensitive to charged interactions.

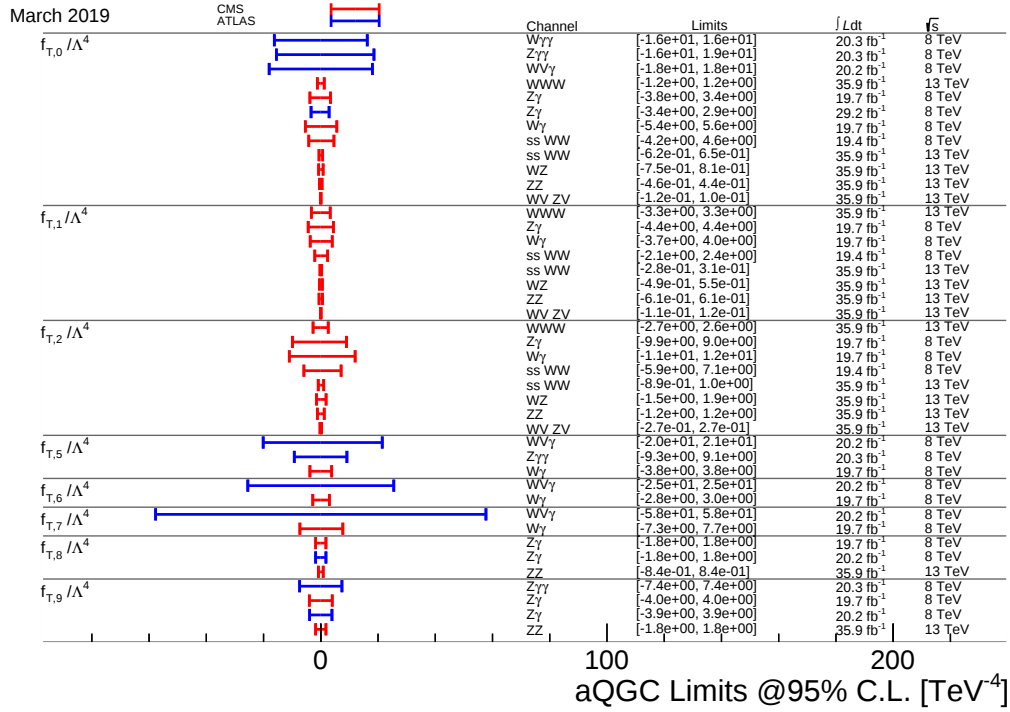


Figure 7.7: Comparison of 95% CL intervals on the T0, T1, and T2 dimension-eight EFT operators obtained by the ATLAS and CMS Collaborations.

7.4 Limits on charged Higgs boson production

The observed event yields by channel, and expected yields adjusted to the best-fit values obtained from a maximum likelihood fit in the charged Higgs boson selection described in Section 6.3 are shown in Fig. 7.9.

Constraints on the production of charged Higgs bosons are derived via a maximum likelihood fit of the expected yields to the data in a control region of events with $m_{jj} > 100 \text{ GeV}$ outside the signal selection and the $m_T(WZ)$ for events the charged Higgs boson selection. The $m_T(WZ)$ distribution for the observed data and expected background, as well as illustrative predictions for charged Higgs boson production in the Georgi-Machacek model are shown in Fig. 7.10. No deviation from the SM prediction is observed, and the resulting observed and expected limits are summarized in Fig. 7.11. Limits are derived on the production cross section in a largely model-

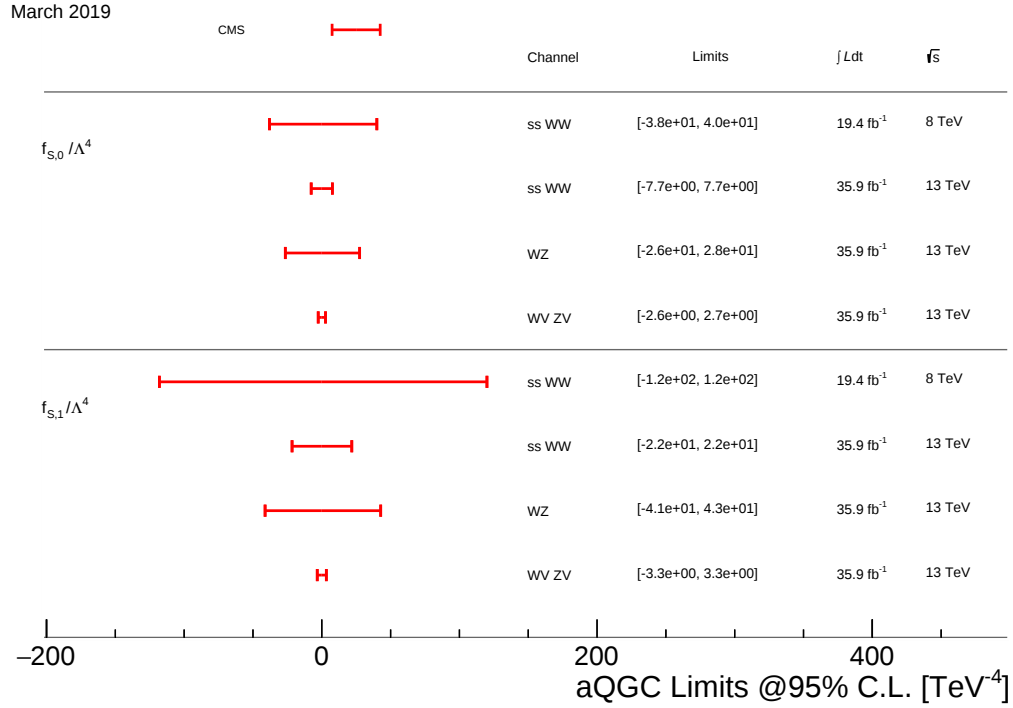


Figure 7.8: Comparison of 95% CL intervals on the S0 and S1 dimension-eight EFT operators obtained by the ATLAS and CMS Collaborations.

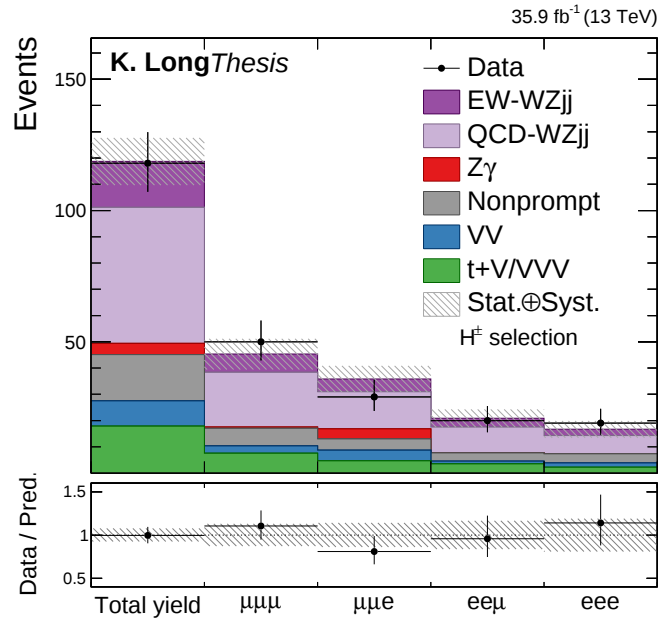


Figure 7.9: Post-fit event yields in the charged Higgs boson search region.

independent way, assuming a narrow-width charged Higgs boson produced via vector boson fusion, and in the Georgi-Machacek model. Results in the Georgi-Machacek model are presented in terms of the parameter s_H and the mass of the charged Higgs boson. These limits are directly comparable to other searches for charged and doubly-charged Higgs bosons in this model.

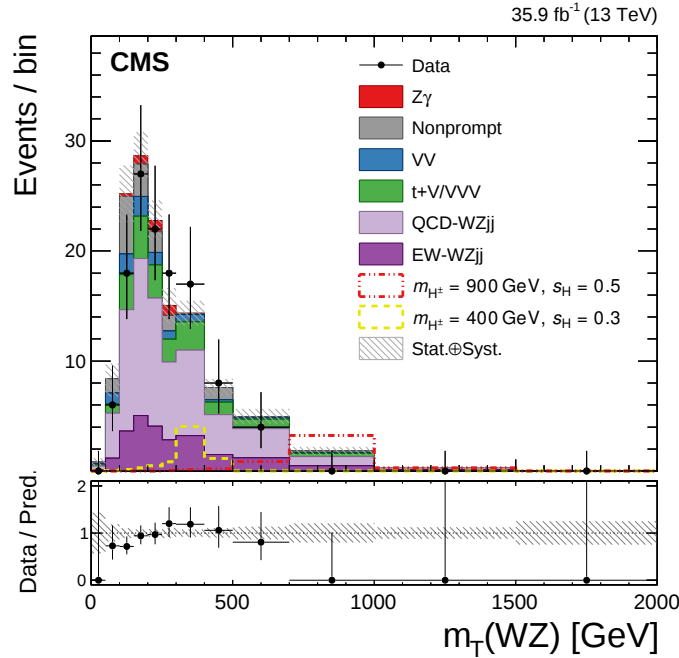


Figure 7.10: The observed $m_T(WZ)$ for events satisfying the Higgs boson selection, used to place constraints on the production of charged Higgs bosons. The last bin contains all events with $m_T(WZ) > 2000$ GeV. The dashed lines show predictions from the Georgi-Machacek model with $m(H^\pm) = 500$ (900) GeV and $s_H = 0.3$ (0.5). The bottom panel shows the ratio of the number of events measured in data to the total number of expected events. The hatched bands represent the total and relative systematic uncertainties on the predicted background yields. The predicted yields are shown with their best-fit normalizations from the background-only fit.

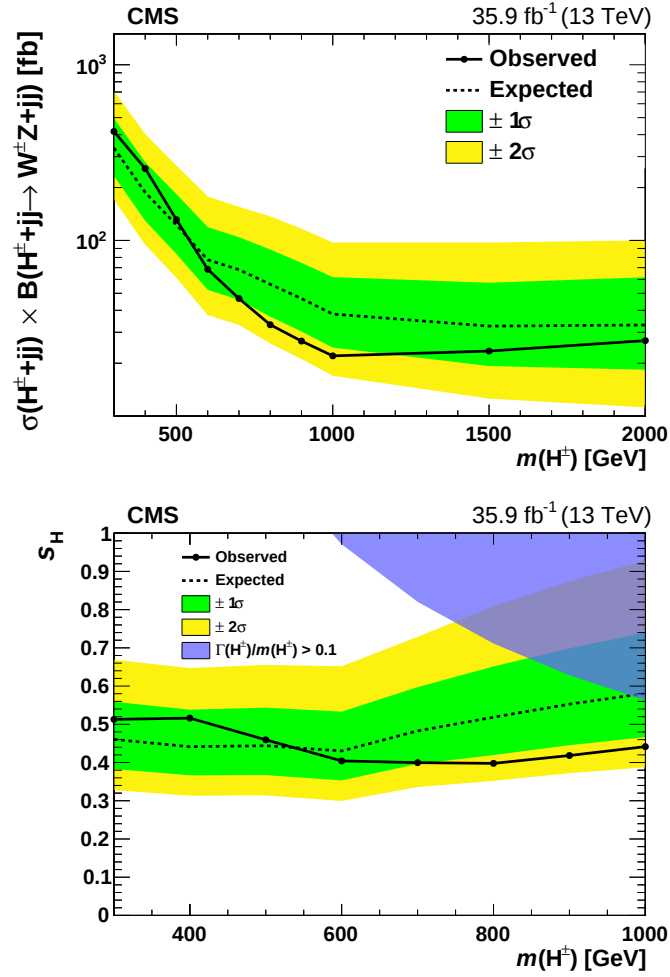


Figure 7.11: Expected (dashed lines) and observed (solid lines) upper limits at 95% CL for the model independent $\sigma(H^\pm) \times \mathcal{B}(H^\pm \rightarrow WZ)$ as a function of $m(H^\pm)$ (left) and s_H in the Georgi–Machacek model (right). The blue shaded area covers the theoretically not allowed parameter space [209].

Chapter 8

Conclusions

8.1 Summary

A measurement of the simultaneous production of a W and a Z boson in association with two jets has been presented, using events where both bosons decay leptonically. Results are based on data corresponding to an integrated luminosity of 35.9 fb^{-1} recorded in proton-proton collisions at $\sqrt{s} = 13 \text{ TeV}$ with the CMS detector at the LHC in 2016. The cross section in a tight fiducial region with enhanced contributions from electroweak (EW) WZ production is $\sigma_{\text{WZjj}}^{\text{fid}} = 3.18_{-0.63}^{+0.71} \text{ fb}$, consistent with the standard model (SM) prediction. The dijet mass and dijet rapidity separation are used to measure the signal strength of EW WZ production with respect to the SM expectation, resulting in $\mu_{\text{EW}} = 0.82_{-0.43}^{+0.51}$. The significance of this result is 2.2 standard deviations with 2.5 standard deviations expected. This is the first study of this process performed by the CMS Collaboration.

Constraints are placed on anomalous quartic gauge couplings in terms of dimension-eight effective field theory operators, and upper limits are given on the production cross section times branching fraction of charged Higgs bosons. The upper limits

on charged Higgs boson production via vector boson fusion with decay to a W and a Z boson extend the results previously published by the CMS Collaboration [103] and are comparable to those of the ATLAS Collaboration [104]. The constraints on dimension-eight effective field theory operators are compatible and complementary to studies performed by the ATLAS and CMS Collaborations using other production and decay channels. For example, the 95% CL limit on the coupling parameters placed on the transverse operator T1 is $f_{T1}/\Lambda^4 \in [-0.49, 0.55]$, compared with the confidence intervals of $[-0.61, 0.61]$ from the ZZ channel [96] and $[-0.28, 0.31]$ in the $W^\pm W^\pm$ channel [94]. These are the first limits for dimension-eight effective field theory operators in the WZ channel at 13 TeV.

8.2 Outlook

The WZjj and EW WZ cross section measurements presented in this thesis are statistically limited. Therefore, a significant reduction in the uncertainty is expected from analyzing a larger data set. The CMS experiment collected data from 13 TeV LHC collisions in 2017 and 2018, corresponding to data sets of 41.1 and 59.7 fb⁻¹ of integrated luminosity. By combining the results presented in this work with an analysis performed on this increased data set, sensitivity to the EW WZ process exceeding 5 standard is expected with only moderate reductions in the systematic uncertainties. Differential measurements of WZjj production in variables sensitive to new physics and corrections from higher perturbative orders will also be possible with larger data sets, and a $\sim 50\%$ reduction in the total cross section uncertainty would be sensitive to the large NLO EW corrections predicted by recent theoretical work [172]. In addition, measurements with reduced uncertainties will definitively determine if the small degree of tension between this result and the result recently

submitted for publication by the ATLAS Collaboration [98] is a result of statistical fluctuations or a consequence of different analysis techniques or interpretations.

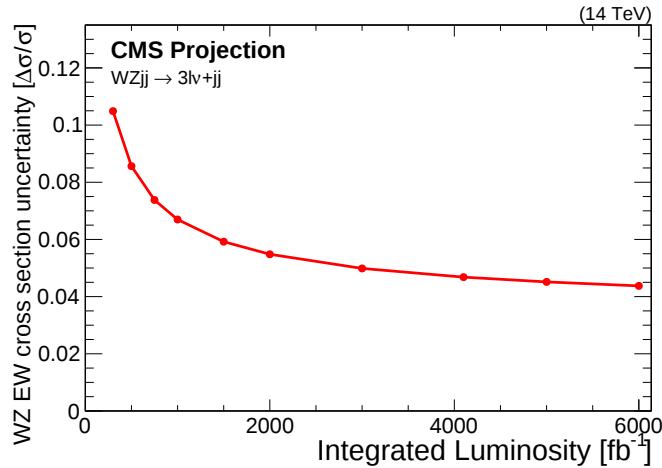


Figure 8.1: Projected reduction in uncertainty for the EW WZ cross section measurement at the HL-LHC. Reproduced from Ref. [210].

The data set collected by the CMS experiment in 2016–2018 constitutes only a small fraction of the data to be delivered by the LHC in its lifetime. Beginning at the end of 2020, the LHC will resume collisions, likely at an increased energy of 14 TeV. A data set corresponding to 300 fb^{-1} is expected to be collected by 2023, at which point the LHC will be upgraded to deliver 5–7 times the current instantaneous luminosities with the high-luminosity LHC (HL-LHC). During more than ten years of data taking the HL-LHC targets 3000 fb^{-1} of data. The analysis presented in this thesis has been projected to the HL-LHC data set by extrapolating the luminosity and estimating the change in reconstruction efficiency due to the harsh pileup conditions at the HL-LHC in Ref. [210]. As shown in Figure 8.1, measurements of the EW WZ process will be possible with percent-level precision. In addition, sensitivity to the longitudinally polarized components of the EW WZ process, which is closely linked to electroweak symmetry breaking and is sensitive to spin-dependent new physics, may be observable by combining high-precision predictions with innovative analysis techniques.

Sensitivity to physics beyond the SM will also improve with the increasing LHC data set. The observable range of production cross sections and operator couplings is expected to improve proportional to the square root of the luminosity, provided the systematic uncertainties remain subdominant. Resonant production of massive particles requires a parton-parton interaction with sufficient energy to produce the resonance, therefore, the observable mass range of charged Higgs bosons or other charged resonances would be dramatically improved by an increase in scattering energy. A high-energy upgrade of the LHC (HE-LHC) that achieves pp collisions at $\sqrt{s} = 27 \text{ TeV}$ by installing hypothetical 16 T magnets in the existing LHC tunnel has been proposed to follow the completion of the HL-LHC project by 2040 [211]. Ambitious proposals for a 100 TeV collider installed in a new 100 km tunnel at CERN or at a new location are also being considered [212].

While it is almost universally believed that new particles or interactions must exist between the electroweak scale set by $\sim m_W \sim m_H$ and the Planck scale $\sim 10^{19} \text{ GeV}$, unlike for the case of the Higgs boson and the LHC [213], a further restriction of this energy scale has not been convincingly established [214]. With the current outlook for the field of particle physics, precision measurements, such as studies of the vector and Higgs boson couplings, remain exceedingly relevant [215, 216]. Small deviations in observables sensitive to new interactions could point to an energy scale of new particles, providing motivation and a target energy regime for a new high-energy collider [217].

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