

Measurements of Z boson production in association with two jets using the ATLAS Run-II dataset

by

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Abstract

The electroweak (EW) production of a Z boson in association with two jets is measured using the full Run-II dataset of the ATLAS experiment. This EW- Zjj process is a fundamental process of the Standard Model (SM), and is sensitive to vector boson fusion (VBF) Z boson production via the WWZ triple gauge vertex. The process is difficult to study, so an advanced methodology is employed to measure the EW- Zjj signal by exploiting topological features. This methodology and the large dataset collected during Run-II have made it possible to measure the cross section of EW- Zjj differentially for the first time. The cross section is measured as a function of four characteristic observables: the invariant mass of the two jet system, the rapidity interval spanned by the two jets, the signed azimuthal angle between the two jets, and the transverse momentum of the Z boson. The observed total fiducial cross section of EW- Zjj is 37.14 ± 3.47 (stat.) ± 5.79 (syst.) fb.

The techniques developed for this analysis can be applied to the measurement of other electroweak processes such as VBF Higgs boson production. EW- Zjj is also an important background for vector boson scattering processes that are of growing interest for searches of deviations from the SM. The differential cross sections themselves provide two avenues for testing the SM. First, the measurements are sufficiently precise as to distinguish between different state-of-the-art theoretical predictions; knowledge gained here is applicable to measurements of other vector boson processes. Second, the differential cross sections are used to test deviations from the SM attributed to higher order corrections in the WWZ vertex by exploiting the sensitivity of a parity-odd observable with an effective field theory approach.

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Chapter 1

Introduction

The success of the *Standard Model of particle physics* (SM) in describing the interactions of fundamental particles is one of the greatest achievements of fundamental physics. However, there are still gaps in our understanding of nature. The SM does not describe the observations of gravitation, dark matter, dark energy or the neutrino masses and mixing. Nevertheless, at current experimentally accessible energies the SM excels in describing the strong and electroweak interactions occurring in particle collisions. In 2012 this was highlighted by the discovery of the Higgs boson by the ATLAS [1] and CMS [2] experiments at the Large Hadron Collider (LHC). This discovery was made with data collected from proton-proton collisions at centre-of-mass energies of 7 and 8 TeV during Run I of the LHC (2009–2013).

Run I was followed by a shutdown to conduct repairs and upgrades of the LHC and the associated detectors, enabling collisions at a higher centre-of-mass energy. Run II saw collisions at 13 TeV and delivered an unprecedented amount of collision data between 2015 and 2018. Analysis of the data collected by the ATLAS and CMS experiments enables a new level of precision measurements of SM predictions. This thesis focuses on one such precision measurement, the electroweak production of

a Z boson in association with two jets (EW- Zjj). The EW- Zjj process is defined by the t -channel exchange of a weak vector boson. As such, this process is sensitive to vector boson fusion (VBF) Z boson production via the WWZ triple gauge vertex and it is of interest in the search for deviations from the SM. The triple gauge couplings were first measured by the experiments at the Large Electron Positron collider (LEP) [3] (1989–2000) that preceded the LHC. The coupling of the WWZ vertex measured with LEP were found to be in agreement with the SM [4]. The EW- Zjj process is difficult to study as the Zjj final state of pp collisions is dominated by Drell-Yan Z production [5] with two jets produced from initial state radiation via the strong interaction. This Drell-Yan process is referred to as strong- Zjj . The nature of the EW- Zjj process leads to a characteristic event topology that differs from that of strong- Zjj .

The cross section of EW- Zjj was first measured by the ATLAS [5] and CMS [6] collaborations with the Run I 8 TeV data. The ATLAS measurement was found to be consistent with the SM predictions of the triple gauge coupling. At the start of Run II the cross section was measured at 13 TeV by the ATLAS [7] and CMS [8] collaborations. The ATLAS measurement used a relatively small dataset collected in 2015 representing 3.2 fb^{-1} and was found to be consistent with the SM predictions. The CMS measurement used the 36 fb^{-1} dataset collected in 2016 and used an effective field theory (EFT) approach to set limits on the anomalous triple gauge coupling, again no deviation from the SM was observed.

With the full 139 fb^{-1} Run II dataset, it is now possible to measure the cross section of EW- Zjj differentially as a function of key observables. The analysis described in this thesis employs an advanced methodology to measure the EW-

Zjj signal by exploiting the topology difference. The techniques developed in overcoming this difficulty can be applied to other electroweak processes similarly dominated, such as the VBF Higgs production. These measurements represent the production rate of EW- Zjj as a function of kinematic properties of the two jet system and of the Z boson itself. Similar differential cross section measurements are also made for the inclusive Zjj process (strong plus EW).

The differential measurements presented in this thesis provide advance our knowledge of fundamental physics along two main avenues: refinement of theoretical modelling and testing the SM by searching for signs of new fundamental interactions. First, the measurements are sufficiently precise as to distinguish between different state-of-the-art theoretical predictions. This help to improve the modelling of the EW- Zjj process as well as related processes such as VBF Higgs production. EW- Zjj is also an important background for Vector Boson Scattering (VBS) processes and diboson and triboson production. VBS is of growing interest at the LHC for searches of deviations from the SM [9]. The EW- Zjj and VBF Higgs signals are both dominated by background process, gluon-gluon fusion Higgs production (with forward two jets) dominates VBF Higgs analagously to the strong- Zjj background of this analysis. Thus techniques developed here to enhance and measure the EW- Zjj signal are likely to be applicable to VBF Higgs. As for the second avenue, the differential cross sections are used to test deviations from the SM attributed to higher order corrections in the WWZ vertex by exploiting the sensitivity of a parity-odd observable with an effective field theory approach.

This thesis is organized as follows. First the necessary theoretical background is outlined in Chapter 2 and the details of the ATLAS detector are described in

Chapter 3. Next, the EW- Zjj analysis itself is presented in three chapters: the technical details of the data, simulation and event reconstruction are in Chapter 4, the analysis strategy and implementation are detailed in Chapter 5 and the measured results are presented in Chapter 6. Chapter 6 also includes the procedure for the correction of detector effects. The results are reinterpreted with an effective field theory interpretation of the Standard Model to study the impact of higher order terms on the WWZ vertex in Chapter 7. Chapter 8 provides a final summary.

Unless a citation is provided in the caption, all figures presented in this thesis were created by myself. I was the main analyser responsible for the work presented in Chapters 4, 5 and 6. The exception to this is Section 6.7, the correction for detector effects, for which Marina Rotaru (IFIN-HH) was responsible. However, I did contribute to this part of the analysis as well. Figures 6.30 and 6.31 were made courtesy of Marina Rotaru. Chapter 7 is the work of Stephen Menary (University of Manchester); it is summarized in this thesis as an important application of the EW- Zjj measurements. The figures included in Chapter 7 were made courtesy of Stephen Menary.

Chapter 2

Theoretical background

Unless otherwise noted, this section is based on References [10–13].

2.1 The Standard Model

The study of high energy particle physics explores three of the four fundamental forces of nature: the *strong*, *electromagnetic* and *weak* forces. The residual strong force binds protons and neutrons together in the nuclei of atoms, the electrons are in turn bound to these nuclei by electromagnetic forces acting on the familiar electric charge, common in everyday life, and the weak force is responsible for the radioactive decay of unstable nuclei. The interactions associated to these three fundamental forces are governed by a single theoretical framework called the Standard Model of particle physics (SM). On a macroscopic level there is another familiar force, *gravity*. However, in the microscopic realm of fundamental particles, gravity is so weak compared to the other three forces that it can be safely ignored.

The SM is a *quantum field theory* that describes the interaction of all known fundamental particles according to the strong, electromagnetic and weak forces.

Each particle is associated to a *field* that carries *symmetry properties* related to the particle's quantum numbers (such as electric charge and spin). These quantum numbers describe internal degrees of freedom of the fundamental particles. A fundamental symmetry of the SM is a *gauge symmetry* related to the invariance of the SM *Lagrangian* under transformations to these internal degrees of freedom. These gauge transformations belong to a mathematical *symmetry group*, that along with the representation of the particle fields, completely determines a *gauge theory*.

The symmetry group that defines the SM is referred to as $SU(3) \otimes SU(2) \otimes U(1)$. It is comprised of unitary ($U(N)$) and special unitary ($SU(N)$) groups of degree N . These groups have a certain number of *generators*, which are transformations that can be repeatedly applied to generate all elements of the group. The SM has twelve non-commuting generators relating to twelve gauged charges, internal degrees of freedom of the particle fields, which take on values according to the eigenvalues of the gauge transformations. $SU(2) \otimes U(1)$ is the group describing both the electromagnetic and weak interactions, collectively called the electroweak (EW) interaction. The familiar electric charge q is the sum of the $U(1)$ generator Y and one of the three $SU(2)$ generators. The strong force is related entirely to the underlying $SU(3)$ symmetry of the so-called *colour charges*.

The particle fields of the SM are each charged under certain generators; they possess quantum numbers that are eigenvalues of the generators (charge operators). The *flavour* (or *species*) of a particle can be described by its quantum numbers, which define how the particle interacts via the three fundamental forces. In the next two sections, the particles of the SM will be introduced, starting with the constituents of matter, namely the spin- $\frac{1}{2}$ fermions, known as *leptons* and *quarks*.

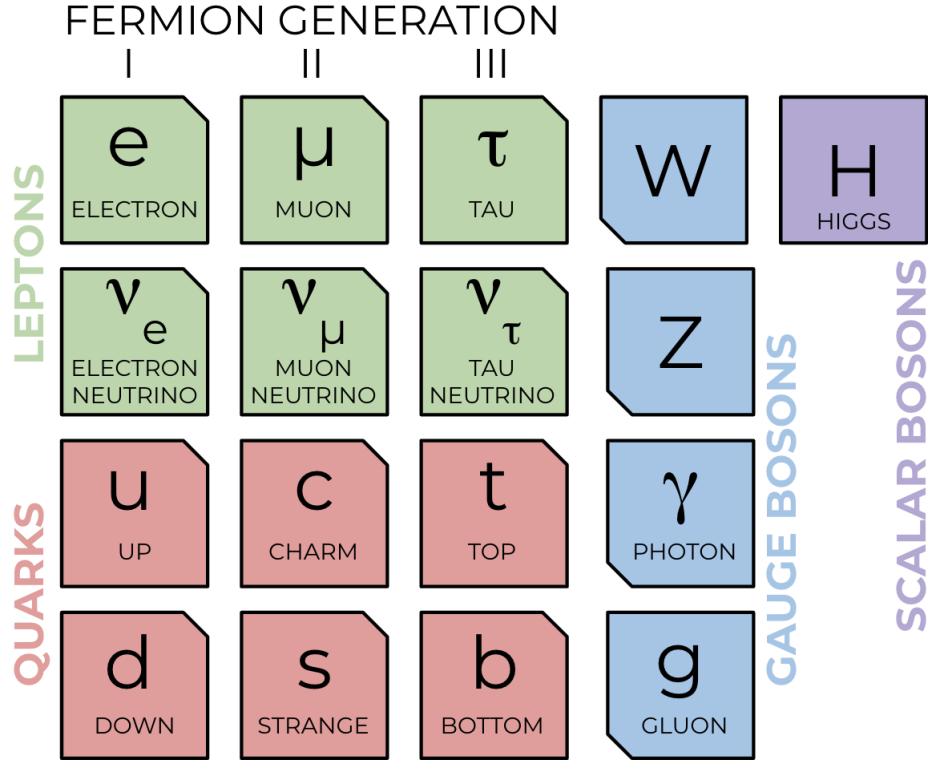


FIGURE 2.1: Particle content of the Standard Model.

Then, the force carrying spin-1 gauge bosons associated to the generators of the SM will be described. Finally, the details of the EW interaction that introduce the scalar Higgs boson and describe the interactions of the weak vector bosons that are the focal point of this physics analysis will be presented.

2.1.1 Fermions

Fermions are spin- $\frac{1}{2}$ particles that bind together to create all matter, the fundamental fermions of the SM consist of two broad categories, the leptons and the quarks. There are six flavours of leptons and of quarks as shown in Figure 2.1. The leptons consist of three electromagnetically charged flavours, the electron e , muon μ and tau

τ particles with $q = -1$, and three neutral particles called neutrinos with flavours associated to each of the charged leptons. The charged leptons differ only in their mass, which leads to the instability of the heavier μ and τ , the electron cannot decay in the SM model as this would violate charge conservation. The neutrinos are known to have mass from neutrino oscillation experiments but this is not included in the SM. Neutrino mass is sufficiently small that they can be treated as massless in the realm of LHC physics.

Quarks also exist in two pairs of three flavours with distinct (fractional) electromagnetic charge, the three up-type quarks: up (u), charm (c) and top (t) have charge $+\frac{2}{3}$, while the three down-type quarks: down (d), strange (s) and bottom (b) have charge $-\frac{1}{3}$. The SM fermions are organized into three generations as depicted in Figure 2.1, the charged leptons and quarks follow an increasing mass hierarchy with increasing generation. Each fermion has a corresponding antiparticle with opposite charge. In the SM, neutrinos exist as both particles and antiparticles, but they could be their own antiparticle. The nature of neutrinos is not yet fully understood.

Quarks differ from leptons in that they are charged under the strong interaction. Each flavour of quark is a *colour triplet*. Quarks are confined to exist only in bound states, never as free particles because of how the interaction strength of the strong force varies with interaction distance and momentum transfer. These bound states, called *hadrons*, are made up of either a quark-antiquark pair (*mesons*) or three quarks (*baryons*). If one was to try and separate a quark bound in a hadron the interaction energy would grow fast enough to create a quark-antiquark pair from the vacuum which will bind to create new hadrons.

In high energy collisions involving hadrons, as momentum is transferred to the

colliding quark inside the hadron it will shatter the hadron creating a spray of new particles, mainly pions, consisting of newly formed quark-antiquark pairs. This phenomenon is called *deep inelastic scattering*.

Fermions also interact via the weak force, and unlike the strong and electromagnetic interactions it does not conserve particle flavour. Leptons and quarks can be converted from one discrete flavour to another via the weak interaction. The SM only contains flavour changing weak interactions that are charged, thus a pair of weakly interacting fermions differ by $\Delta q = \pm 1$. There are no similar flavour changing interactions between like charged fermions, so called flavour changing neutral currents. The strong, weak and electromagnetic forces by which the fermions interact are all mediated by gauge bosons.

2.1.2 Gauge bosons

In Section 2.1, the concept of generators of a symmetry group was introduced. The $SU(3) \otimes SU(2) \otimes U(1)$ symmetry of the SM has twelve generators, and associated to each generator is a spin-1 gauge boson. These bosons act as the force carriers of the three fundamental forces, the 8 generators of the $SU(3)$ colour symmetry are associated to eight gauge bosons called gluons that mediate the strong force. Similarly, the electroweak mediating bosons are associated to the combined $SU(2) \otimes U(1)$ symmetry group. These are the $W^\pm$ and Z bosons and the photon γ . If a gauge symmetry is unbroken like the $SU(3)$ symmetry of the strong force, the mediating bosons are massless like the gluons. Conversely, the EW symmetry is spontaneously broken in the SM leaving only the photon massless and generating relatively large masses for the $W^\pm$ and Z bosons.

The mass m and interaction range r of the force carrying bosons is inversely related through the Heisenberg uncertainty principle $r \sim \hbar/m$, thus the massless photon can carry the electromagnetic force over infinite range. The large mass of the $W^\pm$ and Z bosons limits the range of the weak force to the scale of the atomic nucleus ($\sim 10^{-15}$ m). One might expect the strong force mediated by the massless gluons would have infinite range. However, unlike the photon the gluons carry the colour charge of the strong interaction, which leads to the confinement of both quarks and of gluons themselves.

Thus far, all but one of the particle fields of the SM presented in Figure 2.1 have been introduced. The final category of particles is the scalar boson to which there is a single observed member, the Higgs boson H . The Higgs is related to the mechanism of spontaneous symmetry breaking (SSB) that gives mass to the $W^\pm$ and Z bosons.

2.1.3 Discrete symmetry transformations

There are three important discrete symmetry transformations that can be applied to interactions in the SM: parity (P), charge conjugation (C) and time reversal (T). Observations of these interactions allow physicist to test the conservation or violation of these discrete symmetries.

The parity transformation is the inversion of the three spacial coordinates through the origin, $\mathbf{x} \rightarrow -\mathbf{x}$. When applied to a vector $\vec{v}$, the parity transform produces a vector pointing in the opposite direction, $P(\vec{v}) = -\vec{v}$. There is no evidence of parity violation in the strong and EM interactions. However, in the weak interaction parity is *maximally* violated. A classic example of parity violation is the

process $\pi^- \rightarrow \mu^- + \bar{\nu}_\mu$. The pion is spin-0 and it decays into two spin- $\frac{1}{2}$ particles so in the rest frame of the pion these spins must be oppositely aligned. Observations of this process find that the spin of the muon is always pointing in the same direction as its velocity, called *right-handed* helicity. This implies that the antineutrino is also right-handed. The parity transformed (mirrored) version of this process where the spins are opposite the velocity is not observed.

The charge conjugation transformation converts each particle into its antiparticle. Again, there is no evidence charge conjugation violation in the strong and EM interactions. The combination of applying charge conjugation and parity is the important CP transformation. The CP transformed version of the $\pi^- \rightarrow \mu^- + \bar{\nu}_\mu$ process does indeed occur. This is the similar decay of the positive pion: $\pi^+ \rightarrow \mu^+ + \nu_\mu$. In this case the spins of the antimuon and neutrino will be aligned opposite the velocity, they are *left-handed*. Thus the charge conjugated *and* mirrored (CP transformed) process does occur. However, CP is also not conserved in general by the weak interaction. The CP violation is a very small effect compared to the observed maximal parity violation. It was shown in the decays of the neutral kaons (K^0 and $\bar{K}^0$) that the magnitude of the departure from perfect CP invariance is about 2.3×10^{-3} [14], where 1 represents maximal violation and 0 represents no violation.

Finally, similar to the parity transformation, the time reversal transformation has the effect $t \rightarrow -t$. This transformation converts an interaction to the same process but running in reverse. The combination of CPT is believed to be an exact symmetry of the universe, conserved in all interactions. This implies that if CP is violated then T is also violated to restore the full CPT symmetry.

2.2 Electroweak theory

2.2.1 The Lagrangian and gauge theories

To present the electroweak interaction in detail, the mathematical framework of field theory is needed. The goal of this section is to motivate a unified electroweak theory with massive interacting gauge bosons as in the SM. The equations presented are in natural units where $\hbar = c = 1$.

An important tool is the *Lagrangian* function L , which originates in classical mechanics, and is a function of the generalized coordinates q_i and their time derivatives $\dot{q}_i$. In field theory, the fields ϕ are continuous functions of position and time. The associated coordinates are represented by position time 4-vectors with Greek indices *i.e.* $x^\mu = (t, q_1, q_2, q_3)$. The Lagrangian is a function of the fields and of their time and space derivative which are also represented in a 4-vector, the four-gradient, as

$$\partial_\mu = \left(\frac{\partial}{\partial t}, \frac{\partial}{\partial q_1}, \frac{\partial}{\partial q_2}, \frac{\partial}{\partial q_3} \right). \quad (2.1)$$

The Lagrangian density $\mathcal{L}[\phi, \partial_\mu \phi]$ is typically taken as axiomatic in formulating a field theory, it is related to the classical Lagrangian L by

$$S = \int L dt = \int \mathcal{L}[\phi, \partial_\mu \phi] d^4x, \quad (2.2)$$

where S is the action and the integral d^4x is over the 4 space time coordinates. In the context of field theory the Lagrangian density $\mathcal{L}$ is referred to as simply the Lagrangian. Imposing the condition of local gauge invariance onto the free (non

interacting) Lagrangian for a set of fields gives rise to field theories with interactions. As a first example, consider a spin- $\frac{1}{2}$ field represented by a Dirac spinor ψ with four components satisfying the Dirac equation,

$$(i\gamma^\mu\partial_\mu - m)\psi = 0, \quad (2.3)$$

where m is the mass of the particle associated to the field and γ^μ is a 4-vector of the Dirac matrices. The free Lagrangian for this single field ψ is

$$\mathcal{L} = \bar{\psi}(i\gamma^\mu\partial_\mu - m)\psi, \quad (2.4)$$

where $\bar{\psi}$ is the adjoint (complex conjugate transpose) spinor field of ψ . Imposing a local phase transformation on the field ψ of the form

$$\psi \rightarrow e^{i\theta(x)}\psi, \quad (2.5)$$

where the phase angle θ varies locally as a function of the coordinates x^μ , does not leave the Dirac Lagrangian invariant. Instead it transforms as

$$\mathcal{L} \rightarrow \mathcal{L} + (\partial_\mu\theta)\bar{\psi}\gamma^\mu\psi, \quad (2.6)$$

which motivates the addition of an extra term in the Lagrangian to absorb the extra term in the transformation. Defining the phase angle in terms of a charge q like

$$\lambda(x) = -\frac{\theta(x)}{q}, \quad (2.7)$$

a new Lagrangian can be written as

$$\mathcal{L} = \bar{\psi}(i\gamma^\mu\partial_\mu - m)\psi - q\bar{\psi}\gamma^\mu\psi A_\mu. \quad (2.8)$$

A new vector field A_μ is introduced since it can be imposed to transform under the local phase transformation of ψ like

$$A_\mu \rightarrow A_\mu + \partial_\mu\lambda, \quad (2.9)$$

such that the term is absorbed in the Lagrangian transform. This combination of transformations of ψ and A_μ is called a local gauge transformation, the full gauge invariant Lagrangian also requires a kinetic energy term for the vector field A_μ , which for a spin-1 particle is the Proca Lagrangian

$$\mathcal{L} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} + \frac{1}{2}m_A^2 A^\mu A_\mu. \quad (2.10)$$

where $F^{\mu\nu} = \partial^\mu A^\nu - \partial^\nu A^\mu$ is the electromagnetic field tensor. The second term involving the mass m_A is not gauge invariant, this is resolved by requiring the vector field to be massless. In this case the full gauge invariant Lagrangian is

$$\mathcal{L} = \bar{\psi}(i\gamma^\mu\partial_\mu - m)\psi - \frac{1}{4}F^{\mu\nu}F_{\mu\nu} - q\bar{\psi}\gamma^\mu\psi A_\mu. \quad (2.11)$$

This is exactly the Lagrangian that describes the theory of quantum electrodynamics (QED) with a spin- $\frac{1}{2}$ fermion with electromagnetic charge q and mass m interacting with the massless vector photon field A_μ . Thus the condition of local gauge invariance, when taken as a fundamental symmetry and imposed onto the

free Dirac Lagrangian motivates the introduction of a massless vector field that leads to the description of QED. The Lagrangian can be written more compactly by introducing the covariant derivative $\mathcal{D}_\mu$ to replace ∂_μ where

$$\mathcal{D}_\mu = \partial_\mu - iqA_\mu. \quad (2.12)$$

The QED Lagrangian becomes

$$\mathcal{L}_{\text{QED}} = \bar{\psi}(i\gamma^\mu \mathcal{D}_\mu - m)\psi - \frac{1}{4}F^{\mu\nu}F_{\mu\nu}, \quad (2.13)$$

and with this modification the product $\mathcal{D}_\mu\psi$ transforms exactly as ψ itself transforms, *i.e.*

$$\mathcal{D}_\mu\psi \rightarrow e^{iq\lambda(x)}\mathcal{D}_\mu\psi. \quad (2.14)$$

Thus by replacing the derivative terms $\partial_\mu\psi$ with the proper covariant derivative a gauge invariant Lagrangian is constructed. The other requirement is a locally gauge invariant kinetic energy term for the introduced vector field A_μ , which will be massless. In this example the phase transformation $\psi \rightarrow e^{i\theta}\psi$ that was the basis for the gauge transformation is part of the $U(1)$ symmetry group, this group is simply multiplication of ψ by a 1×1 unitary matrix $U = e^{i\theta}$. This idea can be extended to gauge transformations generated by a variety of symmetry groups, particularly the $SU(2)$ and $SU(3)$ groups of the SM, this more general structure is that of Yang-Mills gauge theory, described in the next section.

In addition to the gauge invariance principle it is also important that the SM is a renormalizable field theory. This property limits the terms in the Lagrangian to be

of at most dimension four (energy⁴) in order to eliminate *ultraviolet divergences*. The combined requirement of gauge invariance and renormalization completely specify the Lagrangian of a field theory for the given gauge group and matter fields.

2.2.2 Yang-Mills gauge theory

This description will use the specific example of an $SU(2)$ gauge symmetry as was first formulated by Yang and Mills [15], but the result can be further generalized to many other group theories [16]. For the $SU(2)$ case consider a theory with two Dirac spinor fields ψ_1 and ψ_2 both with mass m . By organizing the fields in a doublet

$$\psi = \begin{pmatrix} \psi_1 \\ \psi_2 \end{pmatrix}, \quad (2.15)$$

the free Dirac Lagrangian is written

$$\mathcal{L} = i\bar{\psi}\gamma^\mu\partial_\mu\psi - m\bar{\psi}\psi. \quad (2.16)$$

This is a very similar form to the single field case but now the doublet fields transform into one another under an $SU(2)$ local symmetry rotation generated by the Pauli matrices σ^i as

$$\psi \rightarrow \exp \left[i\alpha^i(x) \frac{\sigma^i}{2} \right] \psi. \quad (2.17)$$

In this representation a sum over the repeated i index is implied. This transformation is sufficiently more complicated than the $U(1)$ transformation due to the fact

that there are three non-commuting generators σ_i with commutation relation

$$[\sigma^i, \sigma^j] = i\epsilon^{ijk}\sigma^k, \quad (2.18)$$

where ϵ^{ijk} is the Levi-Civita symbol (the index i here is not to be confused with the imaginary unit). Field theories generated from non-commuting gauge symmetries are called *non-abelian*, and the commuting symmetries like $U(1)$ are associated to *abelian* gauge theories like QED.

The covariant derivative for this $SU(2)$ theory is then

$$\mathcal{D}_\mu = \partial_\mu - igA_\mu^i \frac{\sigma^i}{2}, \quad (2.19)$$

and the product $\mathcal{D}_\mu\psi$ is invariant under the local $SU(2)$ gauge transformation. Here we have introduced a set of three vector fields A_μ^i , one for each of the $SU(2)$ generators and a coupling constant g . As before these fields must be massless to preserve gauge invariance. The gauge invariant kinetic energy term for the fields A_μ^i is constructed from the three field strength tensors $F_{\mu\nu}^i$ which are modified from the abelian case to be

$$F_{\mu\nu}^i = \partial_\mu A_\nu^i - \partial_\nu A_\mu^i + g\epsilon^{ijk}A_\mu^j A_\nu^k. \quad (2.20)$$

With this covariant derivative and these field strength tensors the $SU(2)$ Yang-Mills Lagrangian is written

$$\mathcal{L} = i\bar{\psi}\gamma^\mu\mathcal{D}_\mu\psi - \frac{1}{4}F_{\mu\nu}^i F^{i\mu\nu} - m\bar{\psi}\psi. \quad (2.21)$$

Once again in this Lagrangian there is an implied sum over the index i , this means it will contain terms cubic and quartic in A_μ^i . These terms describe interactions among the vector fields, the SM contains such interactions among the W and Z bosons. However, there is a problem in trying to describe the electroweak interaction with $SU(2)$ Yang-Mills theory and that is the requirement of massless vector bosons for gauge invariance. In formulating this theory, the fermions were organized into a doublet of $SU(2)$ with equal masses and there is no such candidate pair of equal mass fermions in the SM. Both these problems will be addressed by introducing spontaneous symmetry breaking, specifically the Higgs mechanism that is described in Section 2.2.4.

If instead the gauge symmetry transform belongs to $SU(3)$ the Yang-Mills Lagrangian for a colour triplet of equal mass fermions is the basis for the non-abelian gauge theory of the strong interaction, *quantum chromodynamics* (QCD). The eight generators of $SU(3)$ give rise to the eight massless gluons and the Yang-Mills Lagrangian includes terms for the triple and quartic gluon self interactions arising from the non-commuting nature of the generators.

2.2.3 Chirality in the weak interaction

An important feature of the weak interaction of the SM is that it is a chiral gauge theory. The spinor fields ψ can be decomposed in terms of left and right handed chiral eigenstates ψ_L and ψ_R . These are eigenstates of the operator γ^5 , the product of the Dirac matrices, $\gamma^5 = i\gamma^0\gamma^1\gamma^2\gamma^3$. The operators $\frac{(1-\gamma^5)}{2}$ and $\frac{(1+\gamma^5)}{2}$ project the left and right handed states of the spinor field. The charged weak current interaction is

defined by an $SU(2)$ symmetry and couples only to the left handed states of the SM fermions, defined as

$$\psi_L = \frac{(1 - \gamma^5)}{2} \psi, \quad (2.22)$$

and the right handed states are invariant. The weak charged current interaction is flavour changing, it couples together different fermions that differ by one unit of electric charge. Thus in the SM the left handed fermion states are organized into doublets of $SU(2)$ and the right handed states into singlets uncharged under the weak interaction. There are six such left handed doublets, three with a lepton and corresponding neutrino and one for each generation of quarks,

$$\begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L, \begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}_L, \begin{pmatrix} \nu_\tau \\ \tau^- \end{pmatrix}_L, \begin{pmatrix} u \\ d \end{pmatrix}_L, \begin{pmatrix} c \\ s \end{pmatrix}_L, \begin{pmatrix} t \\ b \end{pmatrix}_L. \quad (2.23)$$

Similarly, for antiparticles, the weak interaction couples only to the right handed chiral states, this feature is important in explaining the observed *parity* violation of the weak interaction. The parity transformation (denoted P) inverts the three spacial coordinates $(x, y, z) \rightarrow (-x, -y, -z)$. The famous example of parity violation is the β -decay of ^{60}Co where the emitted electrons are observed to be preferentially emitted in the direction opposite the nuclear spin of the ^{60}Co atom. Applying the parity transform to this picture will lead to one where the electrons are now preferentially emitted parallel to the nuclear spin.

The gauged charge associated to the weak doublets is the so-called *weak isospin*,

I_W , with three components associated to the three generators of $SU(2)$. In the unified electroweak theory the third $SU(2)$ field $A_\mu^{(3)}$ mixes with the $U(1)$ field B_μ to produce the physical Z and γ fields of the SM (this will be described in Section 2.2.5). Associated to the field B_μ is the gauge charge Y , the so-called *weak hypercharge*. The familiar conserved electric charge q is related to the unified electroweak gauge charges as

$$q = I_W^{(3)} + \frac{Y}{2}. \quad (2.24)$$

This relation emphasises the importance of specifically the third component of weak isospin, $I_W^{(3)}$. Typically the term weak isospin refers to the third component. The left handed fermion doublets have $I_W = \frac{1}{2}$, with the raised component having $I_W^{(3)} = +\frac{1}{2}$ and the lowered $I_W^{(3)} = -\frac{1}{2}$.

2.2.4 The Higgs mechanism

Spontaneous symmetry breaking occurs in field theory when the minimum energy state, the vacuum state, of a matter field's potential is not unique. The symmetry of the degenerate vacuum states is spontaneously broken and the system manifests into a single vacuum state. The simplest example to illustrate this is a theory with a single real scalar field ϕ with potential

$$U(\phi) = \frac{1}{2}\mu^2\phi^2 + \frac{1}{4}\lambda\phi^4, \quad (2.25)$$

and Lagrangian

$$\mathcal{L} = \frac{1}{2}(\partial_\mu \phi)(\partial^\mu \phi) - \frac{1}{2}\mu^2 \phi^2 - \frac{1}{4}\lambda \phi^4. \quad (2.26)$$

The formalism of field theory calculations, the Feynman calculus, treats the fields as perturbations about the minimum of the potential. In the example Lagrangian discussed so far the minimum is the trivial one where the fields are zero, however for the case of the potential in Equation (2.25) the minimum is not necessarily at $\phi = 0$. The parameter λ needs to be positive for a finite minimum to exist, but the sign of μ leads to different vacuum configurations as illustrated in Figure 2.2. The case with $\mu^2 > 0$ has a typical single minimum at $\phi = 0$, but the $\mu^2 < 0$ leads to two degenerate minima at

$$\phi = \pm v = \pm \sqrt{\frac{-\mu^2}{\lambda}}. \quad (2.27)$$

In this example the actual vacuum state will be either $\phi = +v$ or $\phi = -v$ breaking the symmetry of the Lagrangian, this v is the so-called *vacuum expectation value* (vev). In either case the Lagrangian can be rewritten in terms of a new field η shifted from the original field ϕ by the constant $\pm v$. This new Lagrangian in terms of η will produce the same dynamics as the original Lagrangian, but has restored the feature of being centred around the minimum of the potential required by perturbation theory. The shifted Lagrangian is written

$$\mathcal{L} = \frac{1}{2}(\partial_\mu \eta)(\partial^\mu \eta) - \lambda v^2 \eta^2 - \lambda v \eta^3 - \frac{1}{4}\lambda \eta^4, \quad (2.28)$$

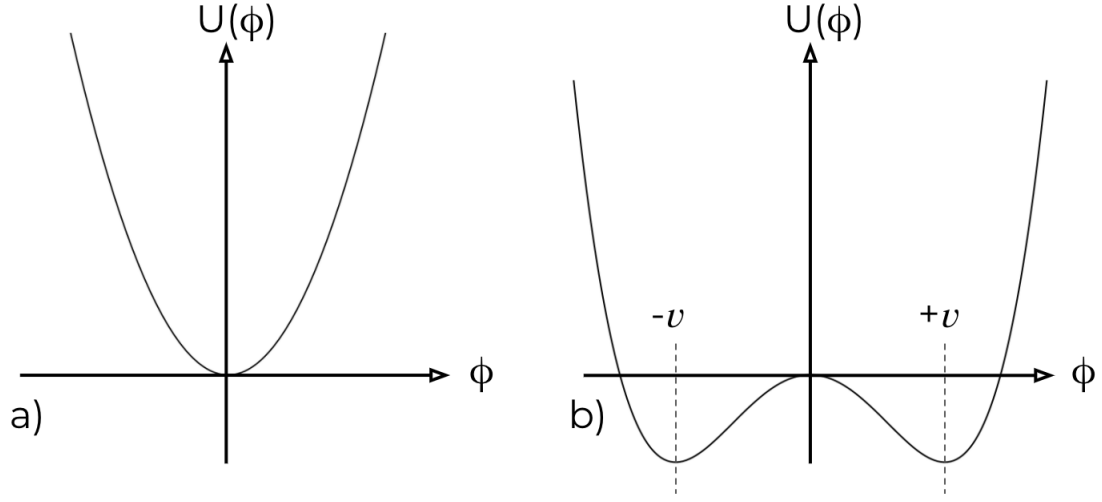


FIGURE 2.2: One dimensional scalar potential with a) $\mu^2 > 0$ and b) $\mu^2 < 0$.

and a mass term for the scalar field can be identified with $m_\eta = \sqrt{2\lambda v^2}$ which is correctly positive unlike the original μ^2 of the ϕ^2 term.

For the next example, if instead the theory contains a complex scalar field expressed in terms of two real scalar fields as $\phi = \frac{1}{\sqrt{2}}(\phi_1 + i\phi_2)$ the Lagrangian is

$$\mathcal{L} = \frac{1}{2}(\partial_\mu \phi_1)(\partial^\mu \phi_1) + \frac{1}{2}(\partial_\mu \phi_2)(\partial^\mu \phi_2) - \frac{1}{2}\mu^2(\phi_1^2 + \phi_2^2) - \frac{1}{4}\lambda(\phi_1^2 + \phi_2^2)^2. \quad (2.29)$$

Similarly, if $\mu^2 < 0$ the potential has degenerate minima. However, this time there is an infinite set of minima defined by

$$\phi_1^2 + \phi_2^2 = \frac{-\mu^2}{\lambda} = v^2, \quad (2.30)$$

as shown in Figure 2.3. Again the Lagrangian is rewritten in terms of the shifted

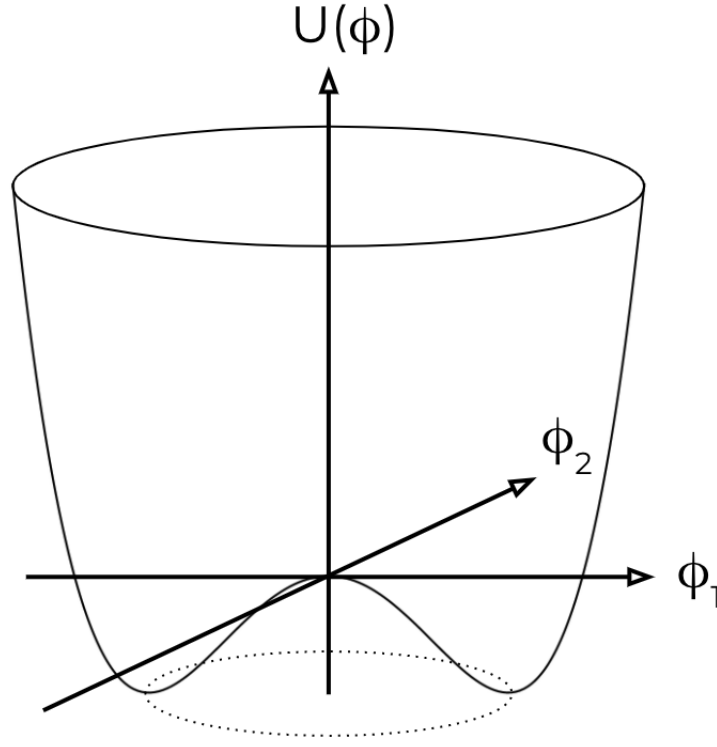


FIGURE 2.3: Complex scalar potential with $\mu^2 < 0$. The dotted circle indicates the set of degenerate minima.

fields centred around a particular minimum, it is convenient to choose the real direction such that $(\phi_1, \phi_2) = (v, 0)$ and define new fields $\phi_1 = \eta + v$ and $\phi_2 = \xi$. For this choice of minimum only the field η will acquire a mass term in the shifted Lagrangian, ξ will be massless. Excitations in the radial direction along the slope of the potential are associated to the massive η field, and ξ corresponds to degenerate rotations around the minimum ring of the potential. This massless scalar ξ is a so-called *Goldstone boson*, these play an important role in the formulation of the Higgs mechanism.

The Higgs mechanism combines spontaneous symmetry breaking and gauge invariance to formulate a gauge theory with massive gauge bosons, like the weak

vector bosons of the SM. An example of this starts from a Lagrangian analogous to the $U(1)$ invariant Yang-Mills Lagrangian but with a complex scalar field ϕ in place of the Dirac field ψ . The Lagrangian is

$$\mathcal{L} = (\mathcal{D}_\mu \phi)(\mathcal{D}^\mu \phi) - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} - \mu^2 \phi^* \phi - \frac{\lambda}{2}(\phi^* \phi)(\phi^* \phi), \quad (2.31)$$

where the covariant derivative contains a gauge boson field A_μ associated to the $U(1)$ symmetry, $\mathcal{D}_\mu = \partial_\mu + igA_\mu$. If $\mu^2 < 0$ the field will acquire a vev as usual with $\phi_1 + i\phi_2 = v$ and the Lagrangian can be rewritten in terms of the new fields η and ξ as

$$\begin{aligned} \mathcal{L} = & \frac{1}{2}(\partial_\mu \eta)(\partial^\mu \eta) - \lambda v^2 \eta^2 + \frac{1}{2}(\partial_\mu \xi)(\partial^\mu \xi) \\ & - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{1}{2}g^2 v^2 A_\mu A^\mu + gvA_\mu(\partial^\mu \xi) - V_{\text{Int}}, \end{aligned} \quad (2.32)$$

where V_{int} contains three and four-point interaction terms between the scalar and gauge fields. In the shifted Lagrangian spontaneous symmetry breaking has introduced a mass term for the gauge boson A_μ related to the vev and to the gauged charge g . Also introduced is a term coupling the gauge boson directly to the Goldstone boson ξ . By a careful choice for the form of the $U(1)$ gauge transformation the field ξ can be eliminated from the Lagrangian, this transformation for A_μ is

$$A_\mu \rightarrow A'_\mu = A_\mu + \frac{1}{gv} \partial_\mu \xi. \quad (2.33)$$

This corresponds to a gauge transformation of the complex field ϕ such that it becomes real valued at every point, this is exactly the shift along the real direction

as before, $\phi = \frac{1}{\sqrt{2}}(v + \eta)$. This is the so-called *unitary gauge*, this choice of gauge removes the degree of freedom associated to the Goldstone boson ξ replacing it with the degree of freedom associated to the longitudinal polarization of the now massive gauge boson. The field η is the Higgs field, usually written as h . In the unitary gauge h and A_μ represent the physical scalar and gauge boson fields. The Lagrangian in terms of these fields is

$$\begin{aligned} \mathcal{L} = & \frac{1}{2}(\partial_\mu h)(\partial^\mu h) - \lambda v^2 h^2 - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{1}{2}g^2 v^2 A_\mu A^\mu \\ & + g^2 v A_\mu A^\mu h + \frac{1}{2}g^2 A_\mu A^\mu h^2 - \lambda v h^3 - \frac{1}{4}\lambda h^4. \end{aligned} \quad (2.34)$$

In this Lagrangian there is a massive scalar Higgs field and a massive gauge boson field both with mass set by the scale of the Higgs vev, $m_A = gv$ and $m_h = \sqrt{2\lambda}v$. Also present are terms for the Higgs-gauge and Higgs self interactions, both features of the SM. To formulate a similar theory that includes the three massive gauge bosons as well as the massless photon of the SM requires starting from a different symmetry group, specifically $SU(2) \times U(1)$.

2.2.5 GWS theory of weak interactions

The model introduced by Glashow, Weinberg and Salam (GWS) gives the experimentally correct description of the weak interaction, for which they were awarded the Nobel prize in 1979. The GWS theory is a spontaneously broken gauge theory that describes the unified electroweak interaction with one massless and three massive gauge bosons. If one were to start from $SU(2)$ gauge symmetry it leads

to a theory with three massive gauge bosons from arguments similar to the spontaneously broken $U(1)$ shown in the previous section. However if instead one starts from a combined $SU(2) \times U(1)$ gauge symmetry with a complex scalar field doublet ϕ the correct SM gauge boson content is reproduced. The theory contains both a neutral scalar field ϕ^0 and a charged field ϕ^+ , this is necessary in the Higgs mechanism to couple the Goldstone bosons to the longitudinal degrees of freedom of the massive neutral Z^0 and charged $W^\pm$ gauge bosons. The scalar field doublet of $SU(2) \times U(1)$ is written

$$\phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix}, \quad (2.35)$$

the symmetry is spontaneously broken and ϕ acquires a vev in the neutral component

$$\langle \phi \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v \end{pmatrix}. \quad (2.36)$$

The covariant derivative in this case introduces gauge fields associated to both the $SU(2)$ and $U(1)$ symmetries,

$$\mathcal{D}_\mu = \partial_\mu - ig_W A_\mu^i \frac{\sigma^i}{2} - ig' \frac{Y}{2} B_\mu, \quad (2.37)$$

where the A_μ^i are a triplet of $SU(2)$ associated to the weak isospin and B_μ associated to $U(1)$ with weak hypercharge Y . The gauge boson mass terms are derived from the term $(\mathcal{D}_\mu \phi)^\dagger (\mathcal{D}^\mu \phi)$ term in the Lagrangian evaluated at the scalar field vev. This

gives the following terms quadratic in the gauge fields,

$$\frac{1}{8}v^2 g_W^2 (A_\mu^{(1)} A^{(1)\mu} + A_\mu^{(2)} A^{(2)\mu}) + \frac{1}{8}v^2 (g + w A_\mu^{(3)} - g' B_\mu) (g_W A^{(3)\mu} - g' B^\mu). \quad (2.38)$$

The first term defines the mass of the massive charged $W^\pm$ bosons defined as

$$W_\mu^\pm = \frac{1}{\sqrt{2}} (A_\mu^{(1)} \mp i A_\mu^{(2)}), \quad \text{with mass} \quad m_W = \frac{1}{2} g_W v. \quad (2.39)$$

The second term when expanded will mix the A_μ and B_μ fields. They can be rewritten in terms of a non-diagonal mass matrix as

$$\frac{1}{8}v^2 \begin{pmatrix} A_\mu^{(3)} & B_\mu \end{pmatrix} \begin{pmatrix} g_W^2 & -g_W g' \\ -g_W g' & g'^2 \end{pmatrix} \begin{pmatrix} A^{(3)\mu} \\ B^\mu \end{pmatrix}. \quad (2.40)$$

The physical Z and γ fields correspond to the diagonal basis for the mass matrix, the physical photon field is denoted A_μ , not to be confused with the triplet A_μ^i . In the diagonal basis this term becomes

$$\frac{1}{8}v^2 \begin{pmatrix} A_\mu & Z_\mu \end{pmatrix} \begin{pmatrix} 0 & 0 \\ 0 & g_W^2 + g'^2 \end{pmatrix} \begin{pmatrix} A^\mu \\ Z^\mu \end{pmatrix}, \quad (2.41)$$

thus the physical fields are the eigenvectors of the mass matrix

$$\begin{aligned} A_\mu &= \frac{g' A_\mu^{(3)} + g_W B_\mu}{\sqrt{g_W^2 + g'^2}}, & \text{with mass } m_A &= 0, \\ Z_\mu &= \frac{g_W A_\mu^{(3)} - g' B_\mu}{\sqrt{g_W^2 + g'^2}}, & \text{with mass } m_Z &= \frac{1}{2}v\sqrt{g_W^2 + g'^2}. \end{aligned} \quad (2.42)$$

These field can be written more neatly in terms of the weak mixing angle $\theta_W = \frac{g'}{g_W}$, such that

$$\begin{pmatrix} A \\ Z \end{pmatrix} = \begin{pmatrix} \cos \theta_W & -\sin \theta_W \\ \sin \theta_W & \cos \theta_W \end{pmatrix} \begin{pmatrix} A^{(3)} \\ B \end{pmatrix}. \quad (2.43)$$

The weak mixing angle provides an elegant relationship between the coupling constants of the weak and electromagnetic interaction and the masses of the massive gauge bosons,

$$e = g \sin \theta_W, \quad (2.44)$$

$$m_W = m_Z \cos \theta_W.$$

The SM Higgs field is introduced here in the same manner as before, defining the field h in the unitary gauge such that

$$\phi = \begin{pmatrix} 0 \\ \frac{v+h}{\sqrt{2}} \end{pmatrix}. \quad (2.45)$$

The mass term in the potential indicates that $m_h = \sqrt{2\lambda}v$ where λ is defined as in Equation (2.31). The term $(\mathcal{D}_\mu \phi)^\dagger (\mathcal{D}^\mu \phi)$ in the Lagrangian includes interaction

terms between the Higgs and the massive gauge fields with couplings proportional to g_W, θ_W and the masses m_Z, m_W ,

$$\begin{aligned} \mathcal{L}[h, W, Z] = & g_W m_W W_\mu^+ W^{-\mu} h + \frac{g^2}{4} W_\mu^+ W^{-\mu} h^2 \\ & + \frac{g m_Z}{2 \cos^2 \theta_W} Z_\mu Z^\mu h + \frac{g^2}{8 \cos^2 \theta_W} Z_\mu Z^\mu h^2 \end{aligned} \quad (2.46)$$

As discussed in Section 2.2.2 the non-abelian nature of the $SU(2)$ theory generates interaction terms between the gauge bosons, of particular interest is the triple gauge vertex V_{WWV} where $V = Z, \gamma$ with couplings

$$g_{WW\gamma} = g_W \sin \theta_W = e, \quad g_{WWZ} = g_W \cos \theta_W, \quad (2.47)$$

as depicted in Figure 2.4. This vertex is probed in the study of electroweak Z plus dijets production and this vertex is sensitive to higher order corrections [17]. So called *anomalous triple gauge couplings* were tested in a previous analysis of the ATLAS Run-1 data [5], where the couplings were found to be in agreement with the SM. In this analysis a different approach was taken to search for higher order effects with an EFT approach that exploits the differential cross section measurements to test the theory. This is detailed in Section 2.6.

The Higgs mechanism is also used to generate mass terms for the fermions, recall that the left handed fermions form $SU(2)$ doublet and the right handed fermions singlets as described in Section 2.2.3. Because of this a mass term like

$$-m\bar{\psi}\psi = -m(\bar{\psi}_R\psi_L + \bar{\psi}_L\psi_R), \quad (2.48)$$

will not be gauge invariant under $SU(2) \times U(1)$. However a term involving the

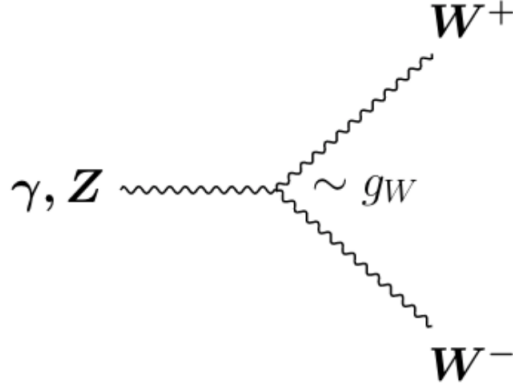


FIGURE 2.4: The triple gauge vertex of the standard model, WWV , where $V = Z, \gamma$.

Higgs doublet will be gauge invariant, with left handed doublet L and right handed singlet R the combinations of $\bar{L}\phi R$ and $\bar{R}\phi^\dagger L$ are invariant. For the example of the electron/electron neutrino doublet and ϕ in unitary gauge from Equation (2.45) there is a term in the Lagrangian of the form

$$\mathcal{L}_{\text{Yukawa}} \supset -\frac{g_e}{\sqrt{2}}(\bar{e}_L e_R + \bar{e}_R e_L) - \frac{g_e}{\sqrt{2}}h(\bar{e}_L e_R + \bar{e}_R e_L), \quad (2.49)$$

where g_e is the *Yukawa coupling* of the electron, a free parameter of the SM. In practice this is chosen to be consistent with the observed electron mass m_e , such that $g_e = \sqrt{2}m_e/v$. Thus the Higgs mechanism introduces a mass terms for the electron and an interaction term between the Higgs and electron fields with a coupling proportional to the electron mass,

$$\mathcal{L}_{\text{Yukawa}} \supset -m_e \bar{e}e - \frac{m_e}{v} \bar{e}e h. \quad (2.50)$$

Mass terms for the up and down type quarks arise in a similar but slightly more complex way but in the end results in one Yukawa coupling for each quark flavour with mass terms and Higgs interaction terms proportional to the quark masses. These Yukawa couplings are all free parameters of the SM constrained by the observed particle masses.

The SM also includes interaction terms between the fermions and gauge bosons, of particular interest is the weak neutral current, the interaction of fermions with the Z boson. For fermion field ψ_f the vertex is

$$V_{\bar{\psi}\psi Z} = \frac{g_W}{2 \cos \theta_W} \bar{\psi}_f \gamma_\mu \left[c_V - c_A \gamma^5 \right] \psi_f Z^\mu. \quad (2.51)$$

Here the couplings c_V and c_A are the *vector* and *axial vector* couplings which are combinations of the couplings c_L and c_R of the Z to right and left handed chiral fermion states. The Z does not couple equally to left and right handed states as is to be expected since the physical field Z_μ is a mixture of $A_\mu^{(3)}$ and B_μ , the couplings are defined

$$\begin{aligned} c_V &= (c_L + c_R) = I_W^{(3)} - 2Q \sin^2 \theta_W, \\ c_A &= I_W^{(3)}. \end{aligned} \quad (2.52)$$

2.2.6 Summary

The Standard Model is a gauge theory containing interacting fields, the observed particles are excitations of these fields consisting of 3 generations of fermions, the quarks and leptons, five force carrying gauge bosons and one scalar Higgs boson. The symmetry group of the standard model is $SU(3) \times SU(2) \times U(1)$, with $SU(3)$

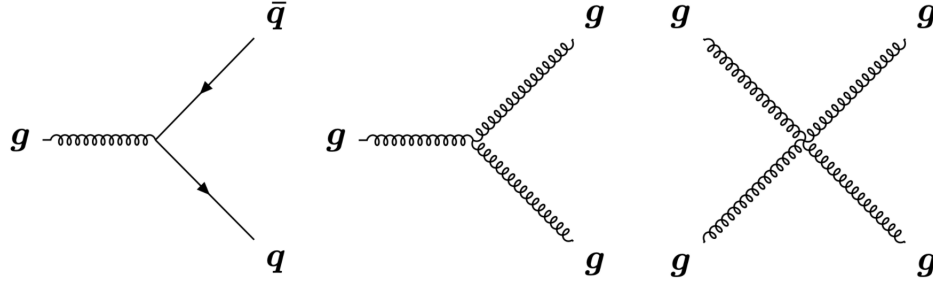


FIGURE 2.5: Feynman diagrams representing the fundamental interactions of the QCD theory of strong interactions, including the quark - gluon vertex and the gluon three and four point self interactions.

corresponding to the colour symmetry of the strong interaction and $SU(2) \times U(1)$ to the unified electroweak interaction. The $SU(2) \times U(1)$ symmetry is spontaneously broken by the Higgs mechanism generating mass terms for the fermions and W/Z bosons. The left handed chiral states of the fermions are organized into doublets of the $SU(2)$ weak interactions, the right handed chiral states are non-interacting singlets. The quarks are colour charged, interacting with the gluon via the strong force and confined to exist in colourless bound states called hadrons, the leptons are not colour charged and the neutrinos are massless in the SM. Particles with electromagnetic charge interact with the photon via the $U(1)$ symmetry with charge $q = I_W^{(3)} + \frac{Y}{2}$ that arises from spontaneous symmetry breaking. Self interactions occur between the weak and gluon gauge bosons due to the non-abelian nature of $SU(2)$ and $SU(3)$. Finally the Higgs field interacts with the massive gauge bosons and the fermions via couplings proportional to particle mass, the particle content of the SM is summarized in Table 2.1. The lowest order (tree level) interaction vertices of the SM are shown for the strong, electroweak and Higgs interactions in Figures 2.5, 2.6 and 2.7, respectively.

Category	Symbol	Mass	Charge (Q)	Spin
Quarks	u	2.16 MeV	$\frac{2}{3}$	$\frac{1}{2}$
	d	4.67 MeV	$-\frac{1}{3}$	$\frac{1}{2}$
	c	1.27 GeV	$\frac{2}{3}$	$\frac{1}{2}$
	s	93 MeV	$-\frac{1}{3}$	$\frac{1}{2}$
	t	173 GeV	$\frac{2}{3}$	$\frac{1}{2}$
	b	4.18 GeV	$-\frac{1}{3}$	$\frac{1}{2}$
Leptons	e	511 keV	-1	$\frac{1}{2}$
	ν_e	~ 0	0	$\frac{1}{2}$
	μ	106 MeV	-1	$\frac{1}{2}$
	ν_μ	~ 0	0	$\frac{1}{2}$
	τ	1.78 GeV	-1	$\frac{1}{2}$
	ν_τ	~ 0	0	$\frac{1}{2}$
Bosons	γ	0	0	1
	g	0	0	1
	$W^\pm$	80.4 GeV	± 1	1
	Z	91.2 GeV	0	1
	h	125 GeV	0	0

TABLE 2.1: Particle content of the Standard Model, masses are quoted from the PDG [18]. Neutrinos are known to have a small but non-zero mass, approximately six orders of magnitude smaller than the electron mass [19].

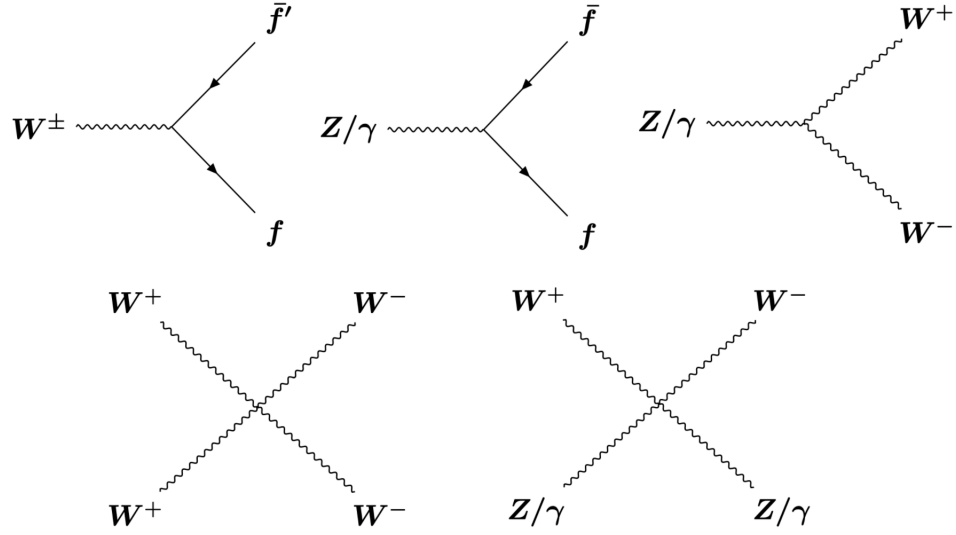


FIGURE 2.6: Feynman diagrams representing the fundamental interactions of the unified electroweak interaction, including the gauge boson - fermion interactions and the gauge boson self interactions.

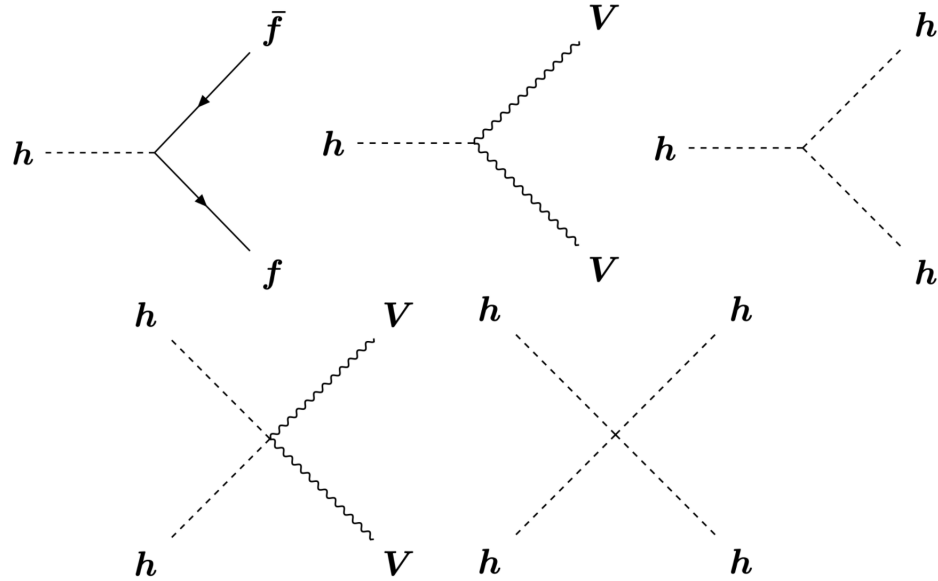


FIGURE 2.7: Feynman diagrams representing the fundamental interactions involving the SM Higgs, including the Higgs-massive fermion Yukawa interaction, the Higgs-gauge interactions where $V = Z, W^\pm$ and the Higgs self interactions.

2.3 Cross sections and pp collision modelling

The analysis presented in this thesis is based on proton-proton collision data at centre-of-mass energy of 13 TeV. When a pair of protons collides at the LHC they are most likely to undergo an *inelastic scattering* where the initial particles are fragmented into a spray of final state hadrons, this process is denoted $pp \rightarrow X$. The rate of occurrence for this process, and for other rarer processes is characterized by the total *cross section* (σ), in this case the total inelastic cross section denoted $\sigma_{\text{inel.}}$.

The driving motivation behind accelerating particles at higher and higher energies is to probe physics at energy scales that have never been studied before. As the collision energy is increased it becomes possible to produce heavier particles as long as the rest mass of the particle is less than the centre-of-mass energy. Collisions producing gauge boson are relatively common at the LHC, but the cross sections of $pp \rightarrow W + X$ and $pp \rightarrow Z/\gamma^* + X$ are $\mathcal{O}(10^6)$ smaller than $\sigma_{\text{inel.}}$, as can be seen in Figure 2.8. This section will address what happens when two high energy protons collide and explain how cross sections are computed as part of pp collision simulation.

Besides the collision energy, another parameter of note is the collision *luminosity*. *Instantaneous luminosity*, L (with units $\text{area}^{-1} \text{ time}^{-1}$), is analogous to the flux of the particle beam, and hence depends on the properties of the beam. The *integrated luminosity*, $\mathcal{L}$, is the time integral of the instantaneous luminosity, this is a measure for the total flux delivered in the beams.

Theoretical predictions of full pp collisions are difficult to compute analytically, so they are simulated statistically using advanced *Monte Carlo* (MC) techniques. The difficulty stems from the composite nature of the protons. The collision is in

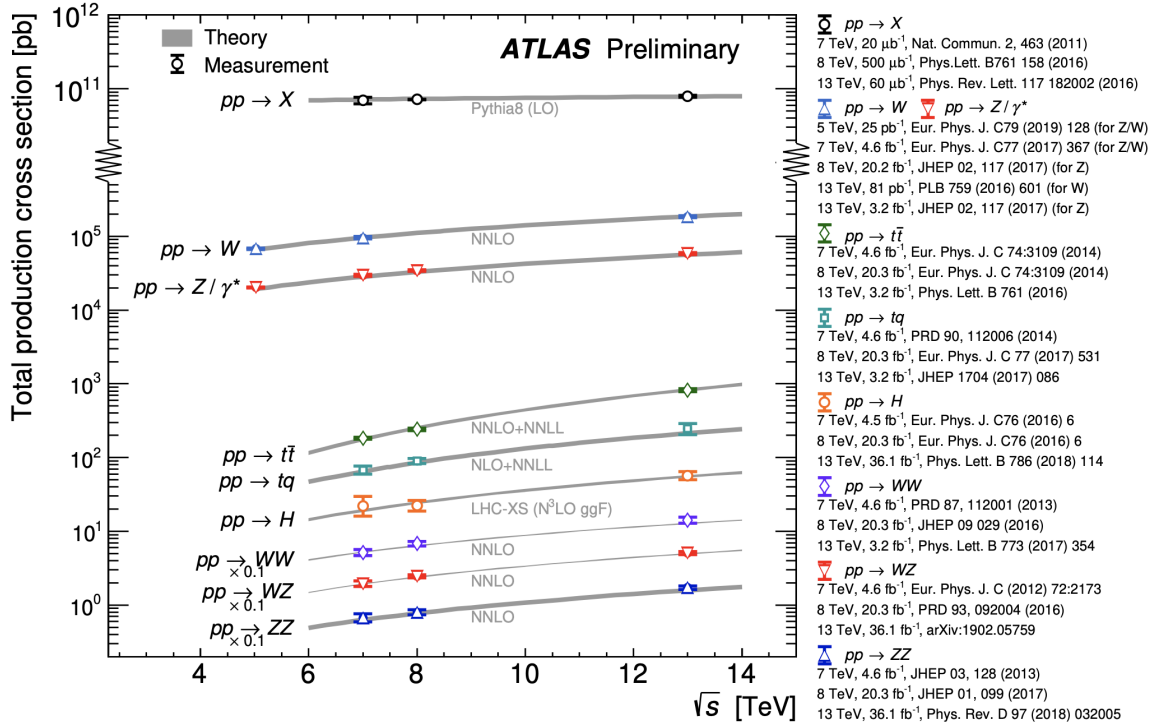


FIGURE 2.8: Summary of total production cross-section measurements by ATLAS presented as a function of centre-of-mass energy from 5 to 13 TeV for a few selected processes. The EW- Zjj process has a cross section similar to that of Higgs boson production ($pp \rightarrow H$) [20].

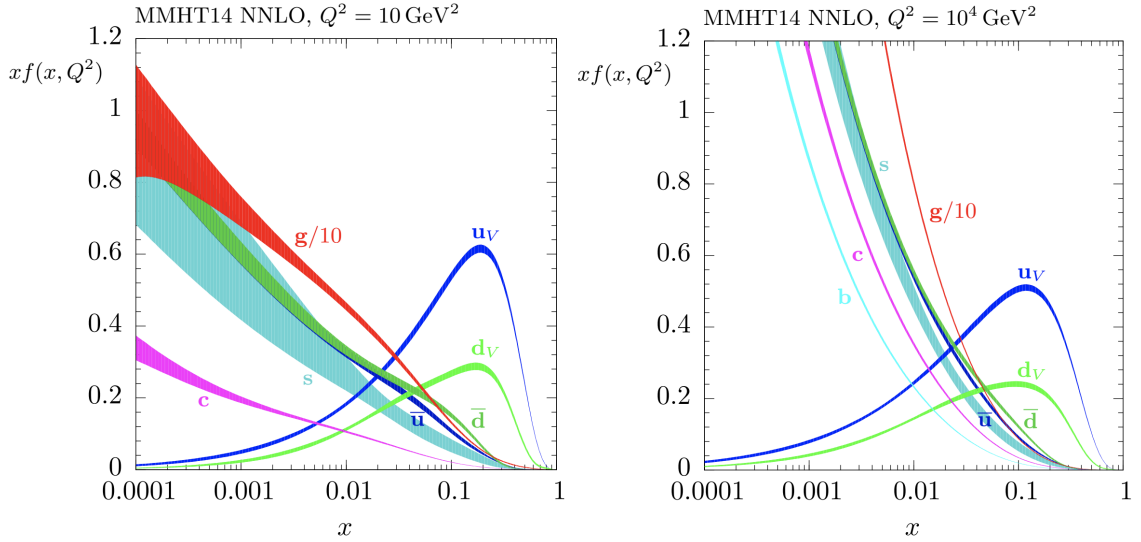


FIGURE 2.9: Proton PDFs at $Q^2 = 10 \text{ GeV}^2$ and $Q^2 = 10 \times 10^4 \text{ GeV}^2$ with associated 68% confidence-level uncertainty bands [21].

fact occurring between constituent *partons* of the protons, carrying some unknown portion x of the collision energy, the so called *momentum fraction*. The colliding partons can be the so-called *valence quarks*, the u and d quarks bound in the proton, or the *sea quarks* that arise from the dynamical interacting system of quarks exchanging gluons inside the proton. The gluons create short lived $q\bar{q}$ sea quarks pairs, thus the interacting parton in a pp collision can be a gluon or antiquark as well. The distribution of momentum fraction for each constituent parton are expressed as *parton distribution functions* (PDFs). PDFs describe the probability of finding a parton with momentum fraction x , displayed typically at a particular value of collision four-momentum transfer squared Q^2 .

In practice, these PDFs are not calculated theoretically. They have a detailed dependence on the dynamics of the proton system and are measured from experiments. A particular set of PDFs are shown in Figure 2.9 for two different values

of Q^2 . The PDFs are used to determine probabilistically which partons interact in a given simulated collision. The process by which the partons interact is called the *hard scatter*. There are several types of hard scatter processes, often classified depending on the produced final state, e.g. QCD *multijet* production, top quark processes, $Z/W^\pm$ boson production or H boson production. The hard scatter processes of interest in a modern collider experiment typically occur at a high $Q^2 > 100 \text{ GeV}^2$, which is within the perturbative realm of QCD where the strong coupling constant is small, $\alpha_s \sim 0.1$, and the quarks of the proton can be treated as free particles.

The hard scatter process can be calculated in a perturbative expansion at a fixed order. The order of the calculations is what defines the accuracy of the simulation in a broad sense, whether it be leading order (LO), next to leading order (NLO), next to next to leading order (NNLO) and so on. In the case of Z boson production the LO process will be the electroweak interaction between the partons to produce the Z boson, LO is sometimes referred to as *tree level* and contains only diagrams with the least number of vertices. Higher orders will add processes with additional strong vertices. At NLO this could be a gluon radiated from an interacting quark, at NNLO this includes virtual correction where a gluon is emitted and reabsorbed (see Figure 2.10). As the order increases the number of processes considered increases rapidly as does the computational complexity, the simulations employed in this analysis were typically of NLO.

The calculation of the scattering amplitude of the hard scatter process, integrating over the full energy spectrum, introduces divergences at the high energy limit (UV) and low energy limit (IR). The UV divergences occur in virtual loops and are eliminated by introducing a *renormalization scale* μ_R that the QCD coupling α_s

becomes a function of. The theory is then defined up to the renormalization scale energy. The most common scheme for absorbing this divergence is the modified minimal subtraction scheme ($\overline{MS}$) [22]. The IR limit introduces divergences when a parton is emitted at a small angle (collinear limit) in the initial state before the hard scatter. This divergence can be solved by factorization, the factorization theorem shows that the cross section can be factorized into a hard part and a normalization from the PDFs. This introduces a *factorization scale* μ_F below which the physics is absorbed into the PDF of the parton. Above the factorization scale the physics is described by the matrix element of the hard scatter process, thus the collinear limit is cut off by integrating starting from μ_F instead of from 0. The cross section for the collision of two hadrons with partons a, b to some final state particles n , $ab \rightarrow n$ can be computed as [23]

$$\sigma = \sum_{a,b} \int_0^1 dx_a dx_b \int d\Phi_n f_a(x_a, \mu_F) f_b(x_b, \mu_F) \frac{1}{2x_a x_b s} |\mathcal{M}_{ab \rightarrow n}|^2(\Phi_n; \mu_F, \mu_R), \quad (2.53)$$

where $f_a(x, \mu_F)$ is the PDF for parton a with momentum fraction x dependent on the factorization scale. Here s is the centre-of-mass energy and $|\mathcal{M}_{ab \rightarrow n}|^2(\Phi_n; \mu_F, \mu_R)$ is the *matrix element* for the process $ab \rightarrow n$ dependent on the final state phase space momenta Φ_n and the renormalization and factorization scales. The matrix element is computed up to some order in $\alpha_s(\mu_R)$ by summing contributions from Feynman diagrams, this defines the precision of the calculation (LO, NLO, NNLO, ...). In LO calculations the dependence on these scales μ_F and μ_R is most significant, at higher order this dependence is reduced with careful choice of scales. The

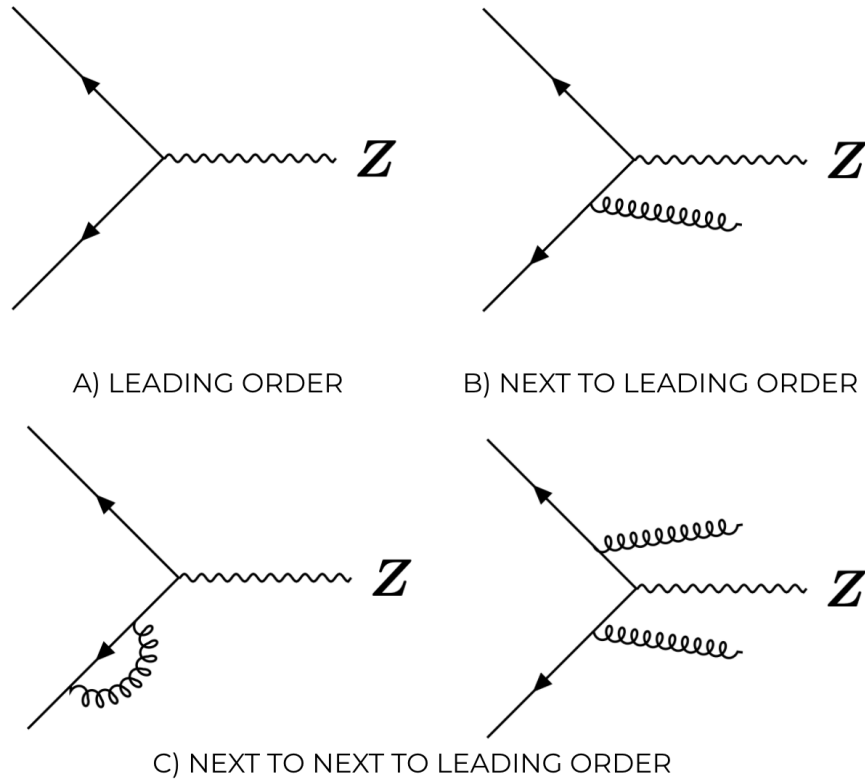


FIGURE 2.10: Depiction of quark-quark fusion production of a Z boson at different orders of perturbative QCD. A) The leading order electroweak process, B) NLO process with emission of a final state gluon, and C) NNLO with a two gluon emissions or a virtual gluon correction on the quark propagator.

uncertainty associated to the choice of μ_F and μ_R is estimated by varying the scales and observing the impact, typically varying up and down by a factor of two.

The calculation of the hard scatter matrix element is the core of the simulation, but to bridge the gap between the hard scatter and the observed detector signal there are several other simulation steps. First is the *parton shower* that simulates how the coloured initial and final state partons split producing additional quarks and gluons. This process is referred to as a shower since it is performed iteratively with

successive splittings from the scale of the hard scatter down to the perturbative limit of QCD around 1 GeV. The parton shower algorithm starts with the LO process and generates multiple splitting sequences with assigned probabilities, thus the LO cross section is partitioned into many cross sections for final states with different distributions and multiplicity of final state quarks and gluons, with sum equal to the cross section of the primary process. Various matching and merging schemes exist to combine the parton shower model with hard splittings that may have occurred in the hard scatter, such as the CKKW* matching procedure [24], with various extensions and improvements [25–27] and the FxFx[†] prescription [28]. Hard splittings are also referred to as *final state radiation* (FSR), typically gluon radiating from a final state quark. When QCD radiation occurs from the partons before the hard scatter it is likewise referred to as *initial state radiation* (ISR).

The parton shower is computed directly via perturbative QCD, at the other end of the shower the short distance (low energy) behaviour of the QCD coupling takes over, α_s is of $\mathcal{O}(1)$ at $Q^2 \sim 1 \text{ GeV}^2$. In this non-perturbative regime the partons become confined into hadrons, this process is called *hadronization*. Hadronization is computed using models tuned to data, the most commonly implemented hadronization models are the Lund string model [29, 30] and the cluster model [31]. The combination of parton shower splitting and hadronization simulates the collimated spray of hadrons observed as so-called *jets* in the detector, with each observed jet

*CKKW is a merging approach based on the computation of all tree-level (LO) matrix elements with up to N partons, where N is as large as the computation can handle. This involves some matching scale below which the parton shower models the collision, which the observables are dependent on, a drawback of using tree-level matrix element.

[†]FxFx is a technique to perform NLO accurate multijet merging. It uses NLO matrix element calculations, which include loop corrections. This has the desired effect of reducing the impact of the matching scale on the observables but the computations are more intensive and the merging more theoretically complex.

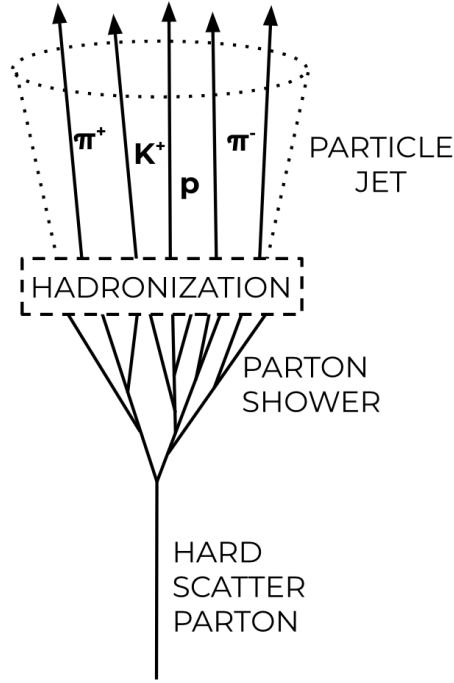


FIGURE 2.11: Diagram depicting the formation of a jet of hadronic particles. For each hard scatter parton the parton shower is simulated with successive splittings down to the perturbative limit of QCD around 1 GeV. Then the hadronization algorithm simulates the formation of a set of collimated hadrons (typically pions, kaons and protons [32]). This is referred to as a *particle jet*.

roughly corresponding to a single parton from the hard scatter process, this process is depicted in Figure 2.11. The most common process observed in high energy pp collisions is the production of a pair of jets. The momenta of the two jets will be perfectly balanced in the centre-of-mass frame of the colliding partons, however in general the incoming partons will not carry the same momentum fraction and thus the two jet pair will appear boosted along the beam direction in the laboratory frame.

The simulation also models and overlays additional activity that does not originate from the primary hard scatter, known as the *underlying event*. This includes the beam remnants and the other partons that did not interact in the hard scatter. These will hadronize, creating radiation preferentially in the *forward* regions, close to the beam line. The remaining partons in the primary hard scatter can also themselves scatter, these are *multiple particle interactions* (MPIs). Most MPI are soft and do not give rise to isolated jets like those coming from the hard scatter. Instead MPIs increase the total amount of scattered energy in the event and increases the number of particles produced in the hadronization stage [23]. Like hadronization the underlying event is simulated from models tuned to data.

A full pp collision is nicely summarized in Figure 2.12, showing the hard scatter, an MPI, initial state radiation, the parton shower and hadronization. The figure depicts an interaction between a quark and gluon parton from the incoming protons, the partons both emit gluons leading to ISR, ultimately the hard scatter interaction depicted is between a pair of gluons. The hard scatter process in this case is $gg \rightarrow t\bar{t}Hg$, where the final state gluon is radiated as ISR. A Feynman diagram for this process is shown in Figure 2.13. The figure depicts three prompt decays, the leftmost line emerging from the hard scatter corresponds to a top quark which emits a gluon (FSR) before decaying ($t \rightarrow W^+b \rightarrow q\bar{q}b$), next is the hard gluon, then the Higgs, which decays via $H \rightarrow b\bar{b}$ and finally the antitop quark which emits two gluons before decaying ($\bar{t} \rightarrow W^- \bar{b} \rightarrow \ell^- \nu \bar{b}$). The outer layer of the figure depicts parton shower level splittings down to ~ 1 GeV where the non-perturbative hadronization model takes over and produces the observable final state hadrons.

The output of such a simulation for an individual event is a list of final state

particles and their four momenta, these are particles with a lifetime τ long enough that they can be detected by an experiment such that $c \tau > 1 \text{ mm}$. Observables of these final state particles from simulation are referred to as being at *particle-level*.

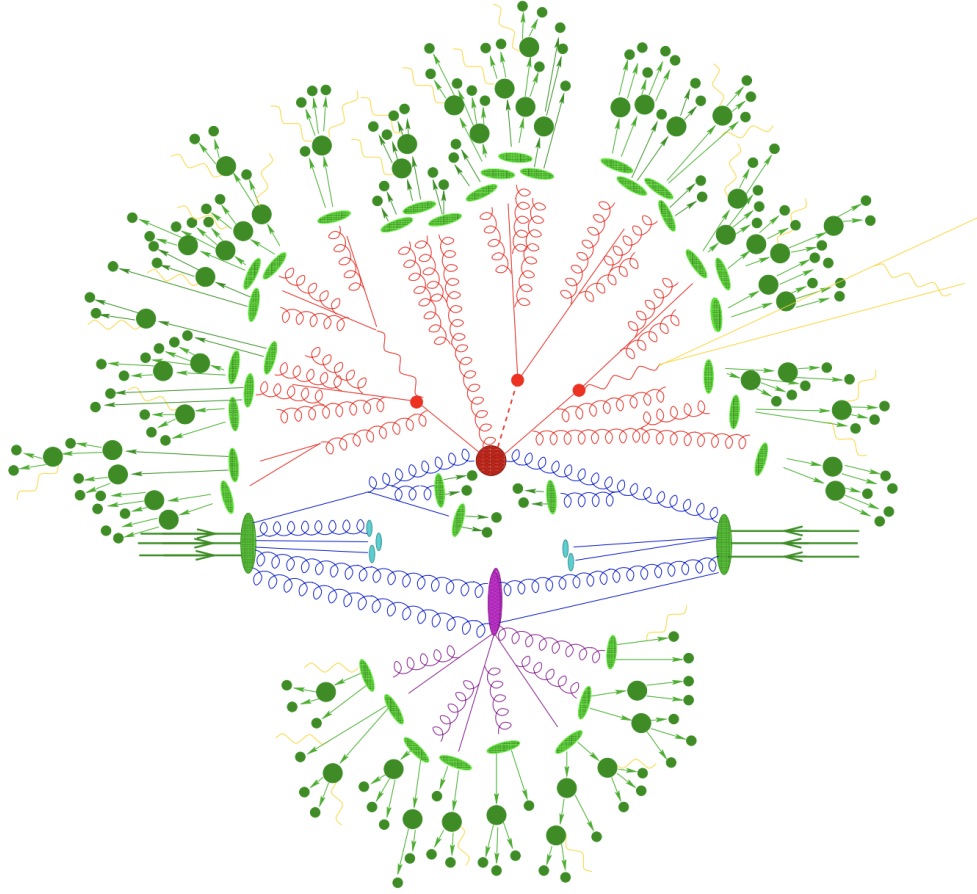


FIGURE 2.12: Pictorial representation of a pp collision. The incoming protons are shown in green with three incoming arrows, corresponding to the three valence quarks, then in the centre the individual partons are in blue. The primary hard scatter is the large red dot, the prompt decay of the hard scatter products are the small red dots and the parton shower are also shown in red. In this case the hard process is $gg \rightarrow t\bar{t}Hg$. A secondary scatter is shown in purple, this is a multiple parton interaction (MPI). The light green ovals around the outside depict hadronization which is followed by hadron decay in dark green. Photon radiation and leptons are shown in yellow [33].

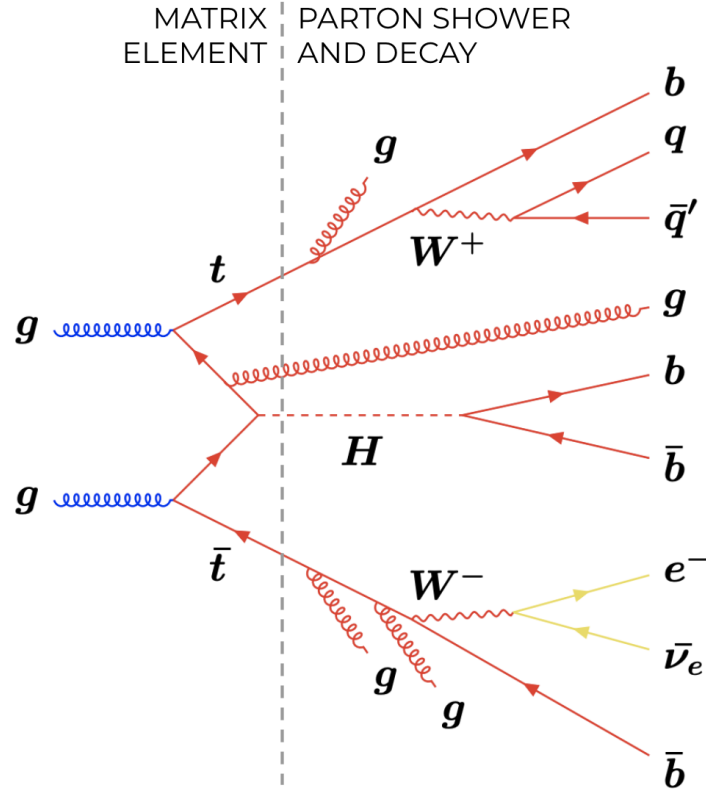


FIGURE 2.13: Possible Feynman diagram of the $gg \rightarrow t\bar{t}Hg$ hard scatter process depicted in Figure 2.12, with internal lines coloured to match. An additional hard gluon is emitted from the internal top quark propagator, a so-called NLO emission. The FSR gluons emitted from the top quarks before their decays are shown as well as the decay products of the top and Higgs prompt decays.

2.4 Zjj production at the LHC

Z boson production at the LHC is most commonly produced by an electroweak process when a quark parton and an antiquark parton from the incoming protons fuse to create a Z boson. When the Z boson decays to a pair of leptons it is called a *Drell-Yan* (DY) process ($q\bar{q} \rightarrow Z \rightarrow \ell^+\ell^-$) as depicted in Figure 2.14. The DY invariant mass spectrum for the dimuon final state is shown in Figure 2.15, the Z boson resonance around 91 GeV causes a large enhancement in the number of observed events.

The Z boson decay modes include pairs of all the fermions of the SM, except the top quark as this would violate conservation of energy since $m_t > m_Z$. The quark decay modes are collectively referred to as *hadronic Z boson decay* and the neutrino modes as *invisible Z boson decay*. The branching ratio to each decay mode is summarized in Table 2.2. The decay modes most often studied by the LHC experiments are $Z \rightarrow e^+e^-$ and $Z \rightarrow \mu^+\mu^-$ as the electrons and muons are well measured in the ATLAS and CMS experiments.

In the DY process, additional partons can be emitted via ISR. Occasionally, these emitted partons can be highly energetic and produce jets. Of particular interest for the study of electroweak Z boson production is the final state with two or more jets $\ell^+\ell^-jj$. The DY process with two or more jets will be referred to as strong- Zjj going forward, diagrams for this process are shown in Figure 2.16. This is the dominant background to the electroweak process and has the same final state. EW- Zjj is defined as all contributions to the $\ell^+\ell^-jj$ final state with a t -channel exchange of an electroweak gauge boson as shown in Figure 2.17.

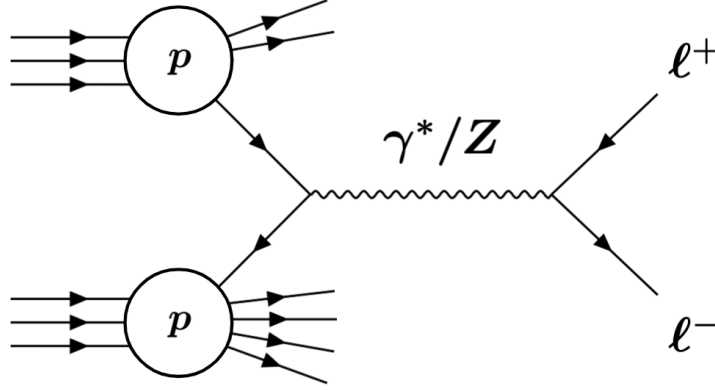


FIGURE 2.14: Diagram depicting a Drell-Yan process via a Z boson in a pp collision.

Z boson decay mode	Branching ratio (%)
e^+e^-	3.3632 ± 0.0042
$\mu^+\mu^-$	3.3662 ± 0.0066
$\tau^+\tau^-$	3.3696 ± 0.0083
invisible	20.000 ± 0.055
hadronic	69.911 ± 0.056

TABLE 2.2: Branching ratios of Z boson decay, hadronic includes decays to quark pairs ($u\bar{u}, d\bar{d}, c\bar{c}, s\bar{s}, b\bar{b}$) and invisible to neutrino pairs of all flavours. Values are quoted from the PDG [18].

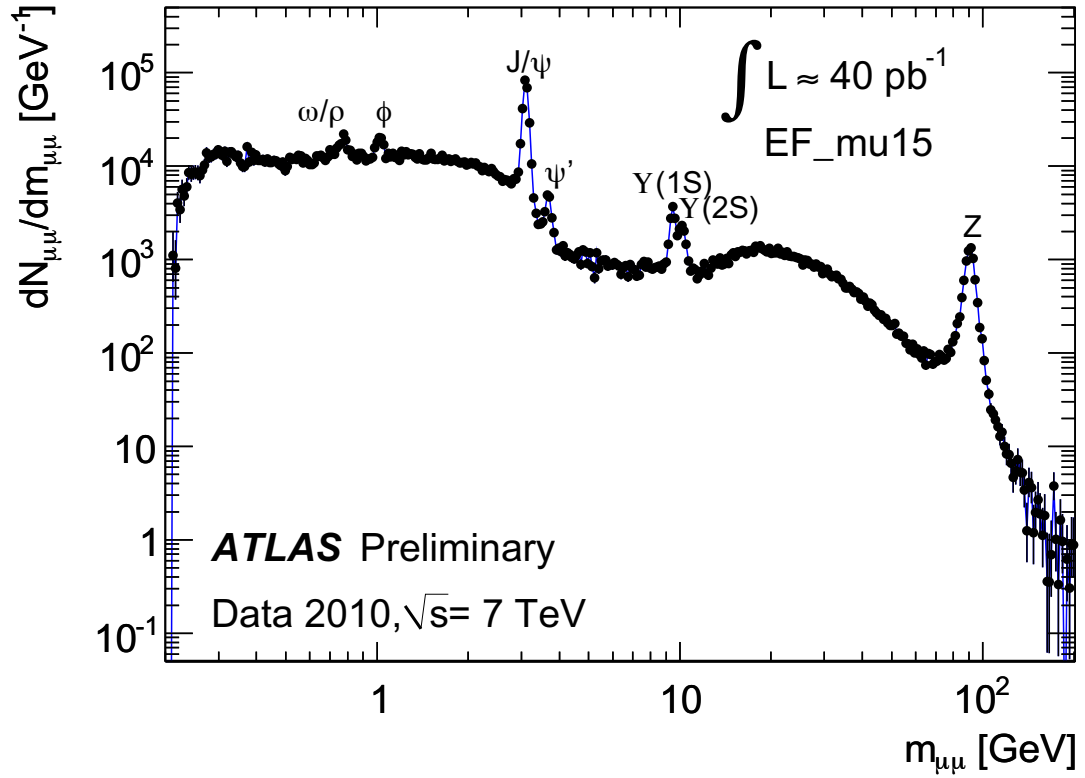


FIGURE 2.15: Di-muon invariant mass spectrum from Drell-Yan events recorded with the ATLAS detector from 2010 collision data at $\sqrt{s} = 7$ TeV. The EF_MU15 label refers to the applied *event filter* trigger, firing on a single $p_T = 15$ GeV muon. This plot corresponds to an integrated luminosity of approximately 40 pb^{-1} [34].

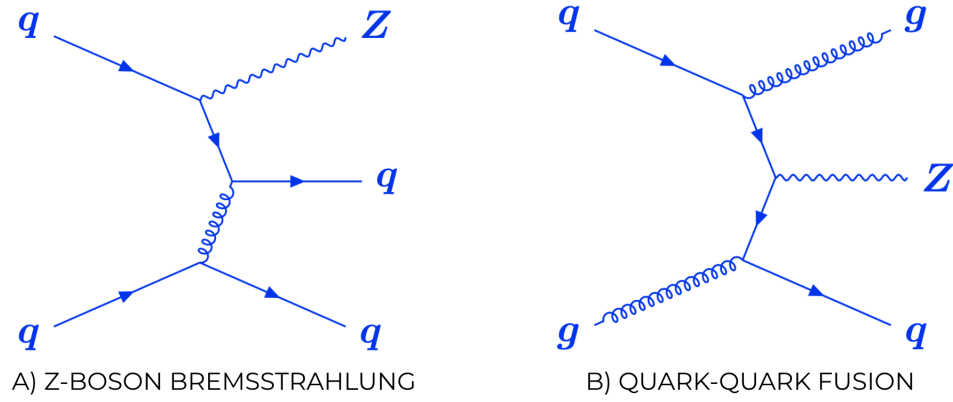


FIGURE 2.16: Representative leading order Feynman diagrams for the strong- Zjj process.

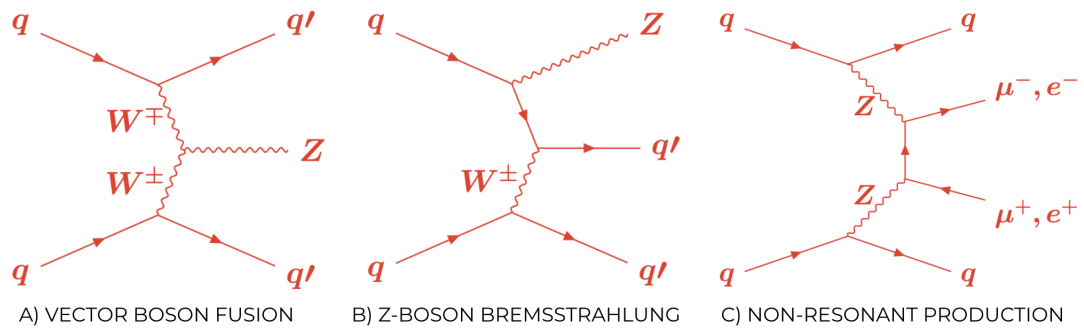


FIGURE 2.17: Representative leading order Feynman diagrams for the EW- Zjj process.

The strong- Zjj and EW- Zjj processes differ in their total production cross sections, the rate of the EW process is $\mathcal{O}(100)$ times lower than that of the strong process. A measured fiducial cross section at 13 TeV of a Z boson produced in association with ≥ 2 jets is (27 ± 3) pb [35], where as another previously measured fiducial cross section of EW- Zjj at 13 TeV is (119 ± 26) fb [7]. These measurements are performed with different fiducial definitions and cannot be directly compared but serve to illustrate the difference in scale between the rates of the two processes. Figure 2.8 also shows large gap between the production rate of $pp \rightarrow Z/\gamma^*$, which is the core of strong- Zjj and the rarer processes such as $pp \rightarrow H$ and diboson production. The total production cross section of EW- Zjj is of the same order of magnitude as $pp \rightarrow H$.

Zjj production also includes a component from diboson production, ZZ or ZW , where one boson decays hadronically and the other to leptons. These processes are not included in the definition of EW- Zjj as there is no t channel exchange of a weak boson as shown Figure 2.18. The topology of these diboson events differs from that of EW- Zjj as the jets originate from the decay of a weak boson instead of from the recoil of the interacting partons from the t -channel exchange.

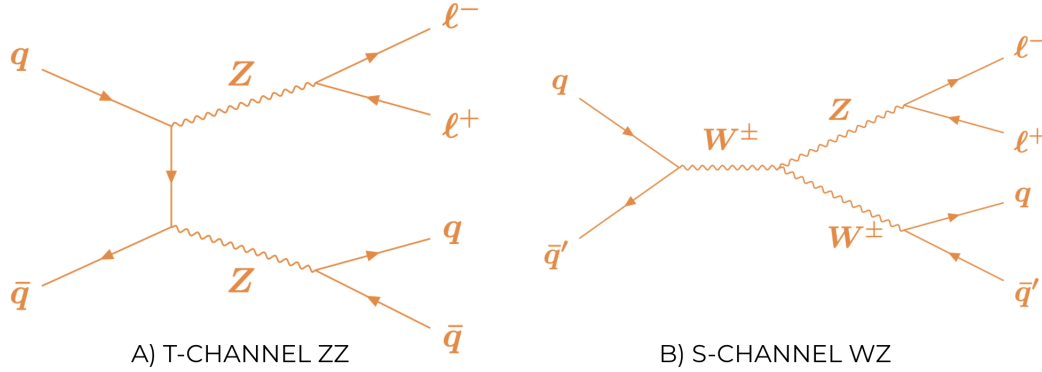


FIGURE 2.18: Representative leading order Feynman diagrams for the diboson Zjj process.

2.5 Zjj event topology

There are several aspects of the event topology of strong- Zjj and EW- Zjj that differ. In this analysis these differences are exploited to enhance the EW- Zjj signal relative to the dominant strong- Zjj background. In order to better present this, the ATLAS coordinate system will first be introduced.

ATLAS uses a right handed coordinate system where the nominal interaction point defines the origin with the beam direction defining the z -axis. The positive x -axis points towards the centre of the LHC ring and the y -axis points straight up. The azimuthal angle ϕ is measured from the positive x -axis increasing towards positive y and the polar angle θ is measured from the beam axis. In general the pseudorapidity η is used in place of θ , defined as $\eta \equiv -\ln \tan(\theta/2)$. For massive objects like particle jets the rapidity y is used instead, defined $y = 0.5 \ln((E + p_z)/(E - p_z))$, where E is the object's energy and p_z its momentum in the z direction. $y = \eta = 0$ in the plane perpendicular to the beam direction and η grows to $\pm\infty$ parallel to the beam. $\eta = y$ in the limit where the particle mass is negligible compared to its total energy. The

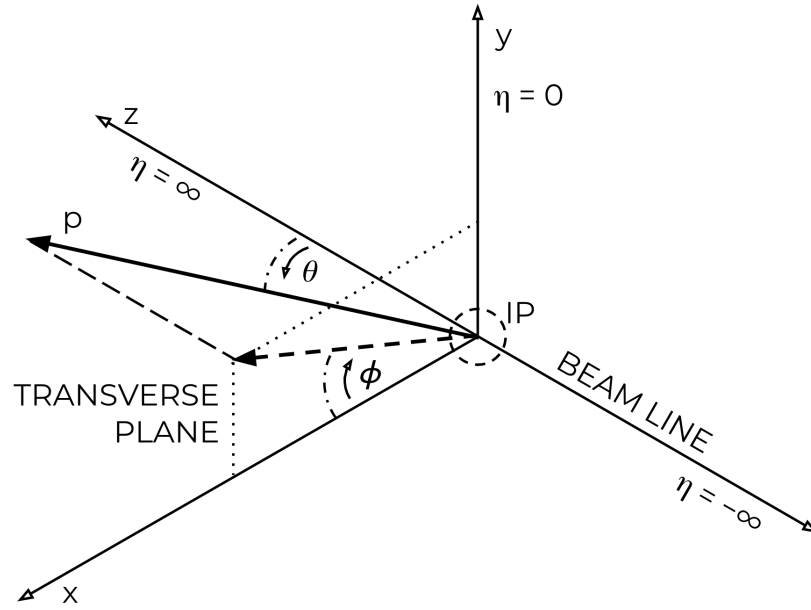


FIGURE 2.19: Coordinate system of the ATLAS detector. A vector p and its projection into the transverse (xy) plane are depicted to illustrate the angles θ and ϕ .

component of a quantity in the $x - y$ plane is referred to as *transverse* with respect to the beam direction, such as transverse momentum p_T of a particular physics object or the missing transverse momentum E_T^{miss} , indicating the presence of undetected particles in an event. Distance in $y - \phi$ space is defined by $\Delta R = \sqrt{\Delta y^2 + \Delta \phi^2}$. The coordinate system is depicted in Figure 2.19.

An important difference between the event topologies of EW- Zjj and strong- Zjj arise due to the so-called *colour connection* between the two interacting quarks in the hard scatter. The topology of the EW- Zjj process can be understood qualitatively from the VBF Feynman diagram of Figure 2.17. The quarks enter with significant boost along the beam axis. After radiating a W boson, the quarks tend to maintain their boost and therefore appear in the forward direction as two back to back jets

in the transverse plane [36]. In strong- Zjj production, the quarks are connected in the diagram by internal quark or gluon propagators as the Z boson is produced via quark-quark fusion. These internal quark/gluon propagators are not present in the purely electroweak process, in EW- Zjj production the incoming quarks interact by a t -channel exchange of a W or Z boson and are thus not colour connected. The Z boson is either produced via vector boson fusion proceeding via the WWZ vertex or it is radiated off of the incoming or outgoing quark. This is an important difference because in these high energy interactions the internal quark/gluon propagators of the colour connected strong- Zjj process will readily radiate one or more quarks or gluons that will appear as additional jets in the final state. Vector boson fusion Z production also requires higher momentum transfer Q^2 than strong- Zjj production, thus the final state objects produced in EW- Zjj will tend to have higher p_T .

The lack of colour connection means that there will tend to be little additional hadronic activity, therefore no additional jets are expected. The rapidity interval between the two *leading* (highest p_T) jets is called the *gap region* and additional jets produced in this gap are called *gap jets*. This region is useful for vetoing the additional hadronic activity expected in strong- Zjj . The lack of additional jets means that the Z boson will tend to balance the dijet system in terms of p_T and that the collision tend to impart more p_T to the Z boson compared to strong- Zjj . The typical EW- Zjj topology is illustrated in Figure 2.20, this will also be referred to as *VBF topology*.

To illustrate the resulting differences in event kinematics, the distributions of a few key observables are presented. First, the two jets produced in EW- Zjj will tend to have more p_T relative to those produced in strong- Zjj . Figure 2.21 shows the jet

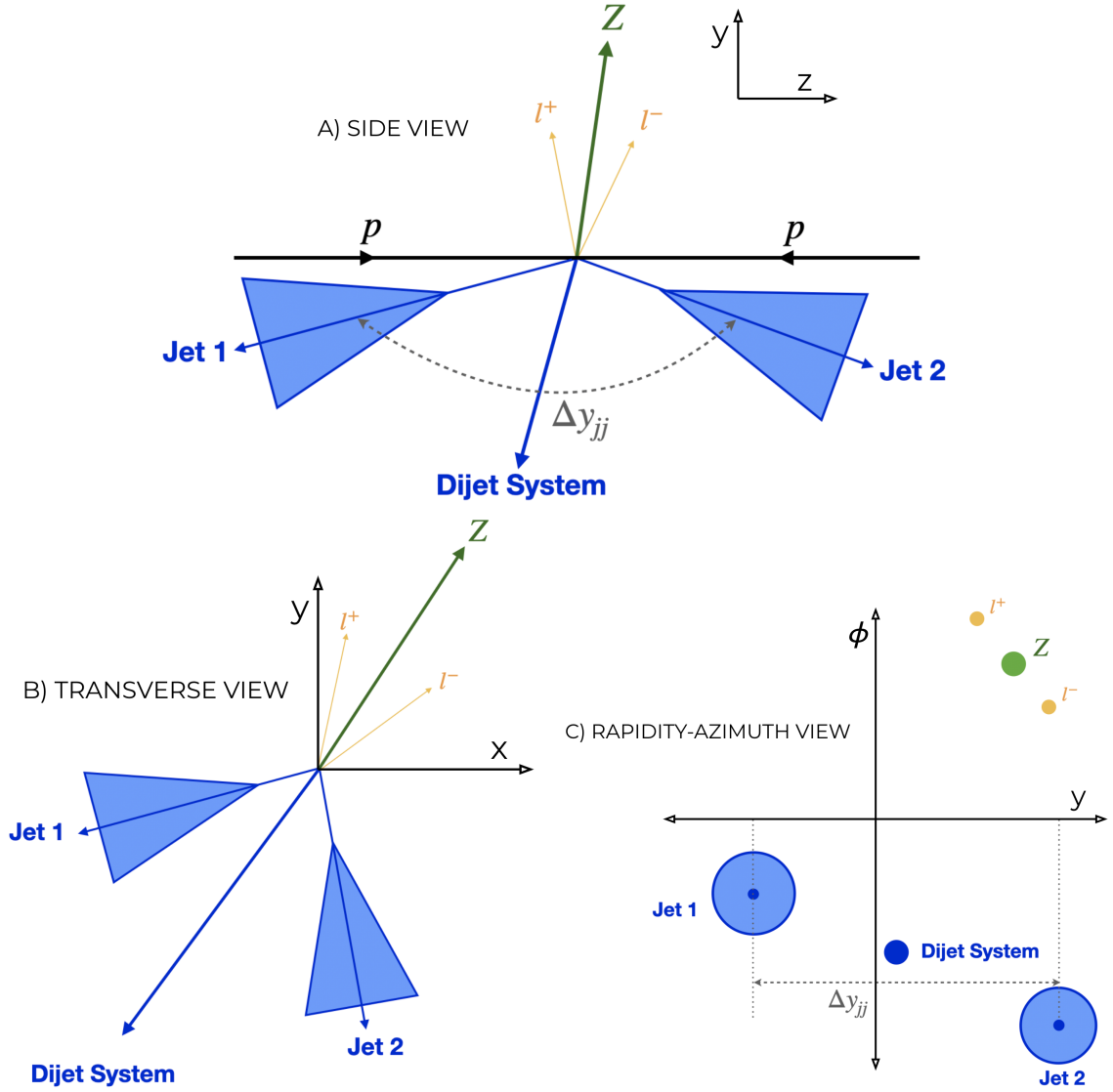


FIGURE 2.20: Diagrams depicting typical EW-Zjj topology from three points of view: A) in the yz plane, B) in the xy plane and C) in the $y\phi$ plane. The dijet system and Z boson are balanced in p_T . The dijet rapidity separation ($|\Delta y_{jj}|$) is labelled to highlight the gap region.

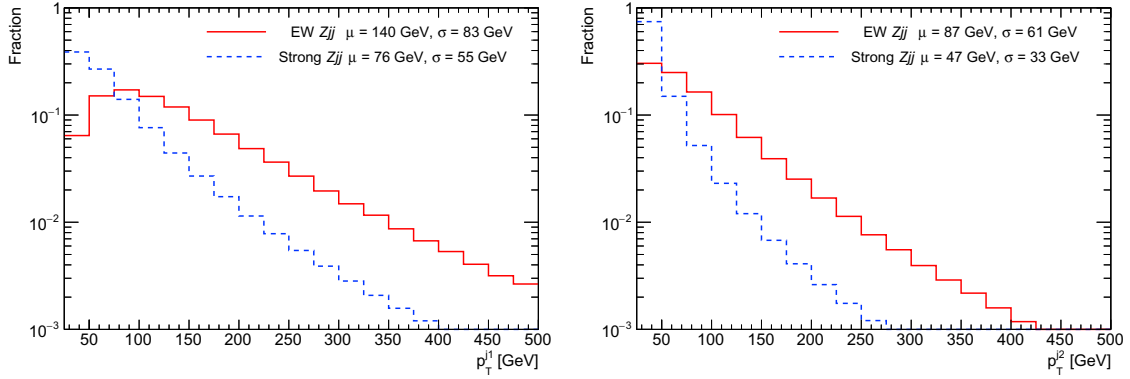


FIGURE 2.21: Normalized particle-level distributions of leading (left) and subleading (right) jet p_T separately for strong- Zjj and EW- Zjj . Events are selected requiring a Z boson candidate and two or more jets with $p_T > 25$ GeV.

p_T spectra of the two leading jets in particle-level simulations. Unless otherwise stated, the jets in the events of this analysis will be ordered by their transverse momentum, *i.e.* $p_T(j1) > p_T(j2)$.

Next, the kinematics differences of the dijet system are captured in two key observables, the dijet invariant mass (m_{jj}) and the dijet rapidity separation ($|\Delta y_{jj}|$). Figure 2.22 shows that the highly correlated m_{jj} and $|\Delta y_{jj}|$ variables tend to be larger in EW- Zjj relative to strong- Zjj . $p_T^{\ell\ell}$ is also larger in EW- Zjj , this is shown in Figure 2.23. Finally, the gap region is examined. The *gap jet multiplicity*, $N_{\text{jets}}^{\text{gap}}$, is defined as the number of jets in the gap region with $p_T > 25$ GeV. EW- Zjj will tend to have lower $N_{\text{jets}}^{\text{gap}}$ compared to strong- Zjj as is also shown in Figure 2.23.

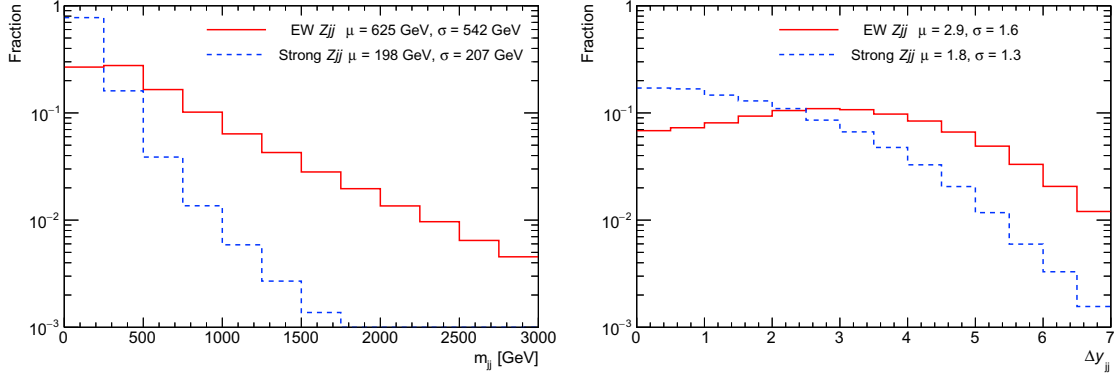


FIGURE 2.22: Normalized particle-level distributions of m_{jj} (left) and $|\Delta y_{jj}|$ (right) separately for strong- Zjj and EW- Zjj . Events are selected requiring a Z boson candidate and two or more jets with $p_T > 25$ GeV.

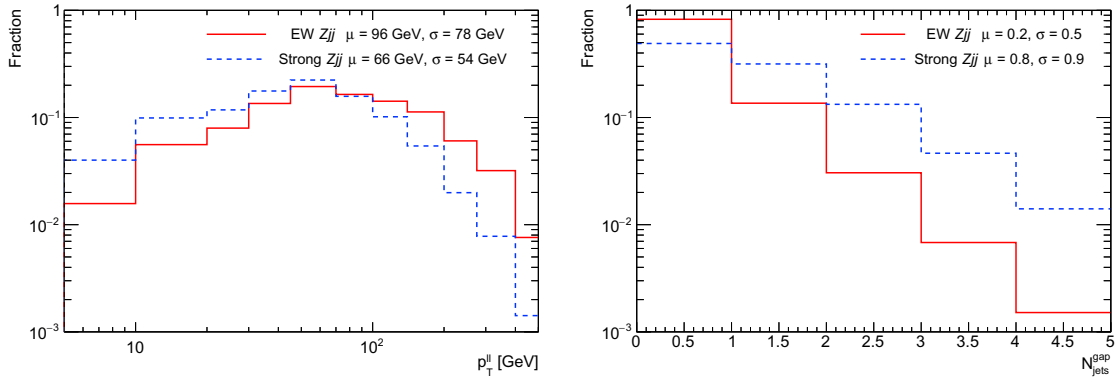


FIGURE 2.23: Normalized particle-level distributions of strong- Zjj and EW- Zjj as a function of $p_T^{\ell\ell}$ (left) and $N_{\text{jets}}^{\text{gap}}$ (right) separately for strong- Zjj and EW- Zjj . Events are selected requiring a Z boson candidate and two or more jets with $p_T > 25$ GeV. The $N_{\text{jets}}^{\text{gap}}$ figure also requires $m_{jj} > 250$ GeV and $|\Delta y_{jj}| > 2.0$.

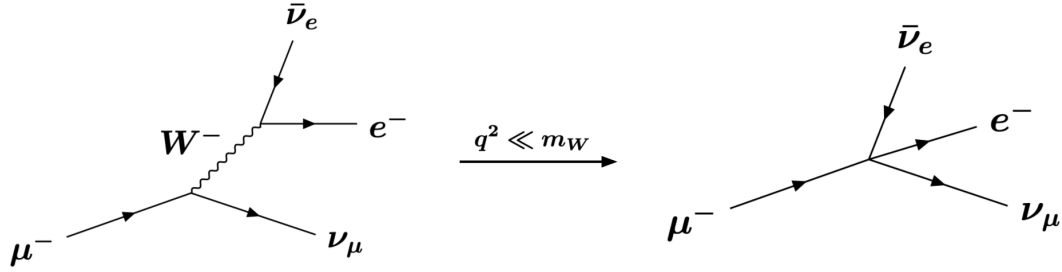


FIGURE 2.24: Feynman diagrams for the W boson mediated muon decay process and the effective field theory approximation reducing the process to a four-point interaction in the low energy regime $q^2 \ll m_W$.

2.6 Incorporation of non-Standard Model interactions using an effective Lagrangian

The classic example of an *effective field theory* (EFT) is the low energy approximation of the weak interaction applicable to the majority of observed particle decays. In these decays the momentum transfer q^2 is much smaller than the mass of the W boson and the interaction can be approximated by an effective four-point interaction without a W boson propagator as depicted in Figure 2.24.

In this example the EFT characterizes the low energy behaviour of a theory characterized by some higher energy scale, in this case the scale is m_W . Another approach is to introduce the effective theory in a model independent way, the *bottom-up* approach. When applied to the SM Lagrangian L_{SM} this approach introduces higher order terms in a systematic way. The renormalizable $L_{\text{SM}}^{(4)}$ contains terms of at most dimension four, in the EFT picture this dimension four part of the SM is assumed to be an effective theory applicable up to some energy scale Λ and the full Lagrangian contains terms of higher dimension suppressed by powers of Λ .

The bottom up approach introduces all terms up to some dimension that satisfy the local $SU(3) \times SU(2) \times U(1)$ symmetry of the SM. In this picture the Lagrangian is written

$$\mathcal{L}_{\text{SM}} = L_{\text{SM}}^{(4)} + \frac{1}{\Lambda} \sum_k c_k^{(5)} \mathcal{O}_k^{(5)} + \frac{1}{\Lambda^2} \sum_k c_k^{(6)} \mathcal{O}_k^{(6)} + \mathcal{O}\left(\frac{1}{\Lambda^3}\right), \quad (2.54)$$

where $\mathcal{O}_k^{(d)}$ represent the higher order terms of dimension d and c_k are the *Wilson coefficients* for each operator. In practice one can enumerate the allowed operators up to dimension six and measure the c_k from observed deviations from the SM. Since all allowed terms have been considered the effective behaviour of the SM is encoded in the c_k no matter what the form of the actual high energy theory is. These higher order terms can be defined in different representations, one such representation is in the *Warsaw basis*, the full list of allowed dimension five and six terms is enumerated in Ref. [37]. As it turns out, the SM gauge symmetry allows for only a single dimension five operator, this operator generates neutrino masses via the *seesaw mechanism*. This is a possible theoretical explanation for the observed neutrino masses not accounted for in the SM. The picture is very different for dimension six where there are 63 allowed operators. These operators modify many aspects of the SM, new vertices appear, SM vertices are modified and there are also shifts to SM parameters like the Higgs vev and weak mixing angle [38].

The higher order terms of interest here are those that modify the WWZ vertex, there are four such operators denoted $\mathcal{O}_W, \tilde{\mathcal{O}}_W, \mathcal{O}_{HWB}$ and $\tilde{\mathcal{O}}_{HWB}$ in the Warsaw basis [37]. This modification has the form $V_{\text{SM}} = a \cdot g + b \cdot g \cdot c_k / \Lambda^2$ where g is the SM coupling and c_k the Wilson coefficient associated to a particular operator $\mathcal{O}_k$. These operators represent new couplings between the Higgs field ϕ and electroweak

fields W_μ^I and B_μ and new self couplings of W_μ^I fields. The form of the operators is shown in Table 2.3. These operators contain the Pauli matrices σ^I and the Levi-Civita symbol ϵ^{IJK} as well as the SM field strength tensors of the electroweak fields, defined as

$$\begin{aligned} W_{\mu\nu}^I &= \partial_\mu W_\nu^I - \partial_\nu W_\mu^I - g\epsilon^{IJK}W_\mu^JW_\nu^K, \\ B_{\mu\nu} &= \partial_\mu B_\nu - \partial_\nu B_\mu. \end{aligned} \quad (2.55)$$

They also contain dual tensors defined as $\tilde{W}_{\mu\nu} = \frac{1}{2}\epsilon_{\mu\nu\rho\sigma}W^{\rho\sigma}$, the two pairs of operators differ only by replacing one $W_{\mu\nu}^I$ tensor with its dual. This has the impact of changing the CP (charge conjugation parity) behaviour of the operator. Those with a dual tensor are CP -odd and those without are CP -even. This is important when considering how the operators will impact differential cross sections measured in parity-even and parity-odd observables. The dimension six operators will impact the scattering amplitude $|\mathcal{M}|^2$ (and therefore the cross section) both with a pure dimension six contribution and via SM - dimension-six interference,

$$|\mathcal{M}|^2 = |\mathcal{M}_{\text{SM}}|^2 + 2\text{Re}(\mathcal{M}_{\text{SM}}^*\mathcal{M}_{\text{d6}}) + |\mathcal{M}_{\text{d6}}|^2. \quad (2.56)$$

The differential cross section of EW- Zjj in the parity-odd observable $\Delta\phi_{jj}^{\text{signed}}$, the signed azimuthal angle between the two jets, is particularly sensitive to the interference term. The details of how this observable is used to set limits on the Wilson coefficients associated to the four dimension six operators is detailed in Chapter 7.

Coupling	Operator
$\mathcal{O}_W$	$\epsilon^{IJK} W_\mu^{I\nu} W_\nu^{J\rho} W_\rho^{K\mu}$
$\tilde{\mathcal{O}}_W$	$\epsilon^{IJK} \tilde{W}_\mu^{I\nu} W_\nu^{J\rho} W_\rho^{K\mu}$
$\mathcal{O}_{HWB}$	$\phi^\dagger \sigma^I \phi W_{\mu\nu} B^{\mu\nu}$
$\tilde{\mathcal{O}}_{HWB}$	$\phi^\dagger \sigma^I \phi \tilde{W}_{\mu\nu} B^{\mu\nu}$

TABLE 2.3: Dimension six operators modifying the WWZ vertex.

Chapter 3

The ATLAS experiment

The ATLAS experiment at the Large Hadron Collider is the result of one of the largest collaborations ever realized in science. The collaboration comprises over 3000 scientific authors from 38 countries spread across the globe. The motivation for the experiment is to improve our understanding of the fundamental constituents of matter. This pursuit has required extreme feats of engineering to construct a particle detector of truly massive proportions and capable of high precision. This chapter will first provide an overview the Large Hadron Collider and then the ATLAS detector itself including the main detector subsystems.

3.1 The Large Hadron Collider

The Large Hadron Collider (LHC) is the world's largest particle accelerator [39], capable of colliding (beams of) protons at centre-of-mass energies up to 14 TeV. The LHC consists of a 27 km main accelerator ring, located 100 m underground straddling the French-Swiss boarder at the European Organization for Nuclear Research (CERN), in Geneva. This main ring is the last step in the accelerator

chain that takes hydrogen gas, uses an electric field to strip the electrons from the hydrogen atoms to produce protons and accelerates these protons to higher and higher energies in stages before finally injecting them into the 27 km ring.

The first stage of the chain is the Linac 2, a linear accelerator that accelerates the protons to 50 MeV. This is the only linear accelerator in the chain, and is followed by a booster ring, and then by two intermediate rings and finally the main ring, all circular accelerators. The booster ring is called the Proton Synchrotron Booster and accelerates the protons from 50 MeV to 1.4 GeV. This is followed by the Proton Synchrotron (PS) itself, which accelerates to 25 GeV, and the Super Proton Synchrotron (SPS), which accelerates to 450 GeV. The SPS injects the 450 GeV protons into the main LHC ring in both the clockwise and counter-clockwise direction via two 2.5 km long injection tunnels [40]. During Run 2 of the LHC (2015 - 2018), the collision energy was 13 TeV, so the main ring accelerated each proton beam from 450 GeV to 6.5 TeV. The full accelerator chain is depicted in Figure 3.1.

The LHC is designed to deliver an instantaneous luminosity of $1.0 \times 10^{34} \text{ cm}^2 \text{ s}^{-1}$. Accounting for the total pp cross section of about 10^{-27} cm^2 (10^8 nb), this works out to roughly 10^9 collisions per second. During Run 2 the LHC delivered 156 fb^{-1} worth of pp collisions in terms of integrated luminosity with a peak instantaneous luminosity reaching $2.1 \times 10^{34} \text{ cm}^2 \text{ s}^{-1}$. These collisions take place at four *interaction points* (IPs) around the LHC ring as shown in Figure 3.1. Located at each interaction point is a main experiment, namely ATLAS [1], CMS [2], LHCb [42] and ALICE [43]. The protons circulated in the LHC are packed into bunches of $\mathcal{O}(10 \text{ cm})$ in length [40], each containing $\mathcal{O}(10^{11})$ protons. In Run II these bunches were spaced 25 ns apart, equivalent to 7.5 m for an object travelling at the speed

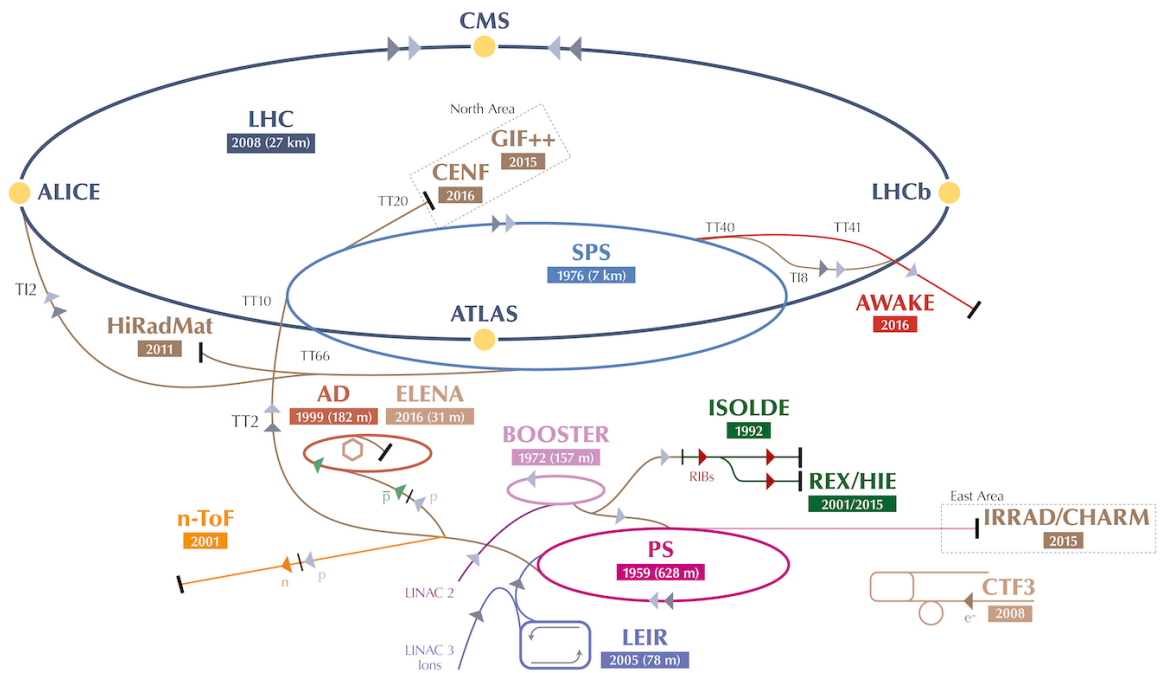


FIGURE 3.1: The CERN accelerator complex, the LHC chain is Linac 2 - Booster - PS - SPS - LHC. The interaction points are shown as yellow circles on the LHC main ring [41].

of light, therefore there is room for ~ 3600 bunches around the 27 km LHC ring. However, in practice not every bunch slot is filled with protons, the maximum number of circulating bunches was 2556 during Run II [44].

The charged protons are bent around the ring via superconducting dipole magnets generating a bending field perpendicular to the proton velocity. The LHC accommodates 1232 dipole bending magnets cooled to temperatures below 2 K and generating fields in excess of 8 T [40] located in eight 2.45 km long curved sections of the ring. Between the curved sections are eight 545 m long straight section containing various types of magnets designed to control the particle beam and to facilitate collisions at the interaction points. Around the IPs quadrupole magnets are used to focus the beams into a radius of $8.5 \mu\text{m}$ [18] to increase the luminosity of the collisions. The beams are collided not head-on but at a small crossing angle of $\mathcal{O}(100 \mu\text{rad})$. This is to avoid undesired collisions between neighbouring bunches as they briefly travel in a single beam pipe. This angle is precisely controlled and is changed over the course of an individual *run* of data taking. A run consists of collisions with two continuously circulating protons beams that lasts around 10 hours, as the beams circulate there is a rate of proton loss that will reduce the collision luminosity over the course of the run, this is somewhat counteracted by reducing the crossing angle of the beams in stages as the run progresses. At the end of a run the beam is discarded (*dumped*) and there is a downtime of a few hours after which a new beam is injected and a new run begins.

The two general purpose particle detectors are the ATLAS (A Toroidal LHC ApparatuS) and CMS (Compact Muon Solenoid) experiments located opposite one another at roughly the southern and northern points of the LHC, respectively.

There is a large overlap in the physics program between these two independent collaborations in order to provide a source of verification for the measurements performed. The ATLAS experiment will be described in detail in Section 3.2. In general the ATLAS and CMS experiments are designed to capture sufficient information about the final state particles produced in the pp collisions to recreate the entire event, thus the detectors cover nearly the full 4π of solid angle surrounding their interaction points. The experiments are both designed with layers of different detectors including an *inner tracker*, *calorimeters* and a dedicated *muon spectrometer*. Both experiments also employ magnetic fields to curve charge particles for momentum determination, CMS uses a single solenoid field and ATLAS a inner solenoid field and outer toroidal field. The physics programs of both experiments include SM precision measurements, Higgs boson discovery and subsequent Higgs sector studies and various beyond the Standard Model searches.

The LHCb (**LHC beauty**) experiment is located at the south-eastern IP of the LHC. It features an asymmetric design that excels at identifying the decay of b quark containing mesons. The LHCb physics program includes measurements of rare decay channels, CP violation and searches for flavour changing neutral currents. Located at the south-western IP is the ALICE (**A Large Ion Collider Experiment**) experiment, a general purpose particle detector designed primarily to study lead ion collisions (Pb-Pb) provided in dedicated runs of the LHC at a centre-of-mass energy of up to 5 TeV. The physics program of ALICE focuses on the study of quark-gluon plasma, a state of matter in which quarks and gluons are no longer confined produced due to the extreme energy density provided in nucleon-nucleon collisions.

3.2 The ATLAS detector

The ATLAS detector is designed to measure the particles produced in LHC collisions. When a pair of bunches from the two proton beams cross one another, many pp interactions can simultaneously occur in this *bunch crossing*. Collisions occur at the *beam spot*, the region where the beams cross. The position and size of the beam spot are determined by the operating parameters of the LHC. The detector records *events* consisting of detector read outs corresponding to a particular bunch crossing. In order to reconstruct the outgoing particles the detector is comprised of layers of detector subsystems optimized for different varieties of particles. To successfully measure the physical processes of interest each subsystem must cover a sufficient area of acceptance and perform at the required resolution. The descriptions presented in this section are based primarily on information from Ref. [1].

Besides the detector performance at the level of individual particle detection, the experiment as a whole must be able to make a fast decision as to whether a given event is of interest and should be recorded, this is handled by the *trigger system*. In Run II data taking the rate of collisions (proton *bunch crossings*) is approximately 30 MHz, the hardware based Level-1 (L1) trigger reduces this rate down to 100 kHz and the software based high level trigger (HLT) reduces even further to approximately 1 kHz [45]. The performance of the trigger system is crucial in order to keep the amount of recorded data to a manageable level while maintaining a high efficiency of recording the desired events that potentially contain the physical processes of interest, such as the production of Z bosons.

The L1 trigger is implemented in fast, custom-made electronics. This trigger

system combines reduced-granularity information from the calorimeter (L1Calo) and muon system (L1Muon) to make trigger decisions to identify signals that could be of interest for physics analyses, including the regions of interest (RoIs) in the detector. These are passed for further analysis by the HLT. For Run-II a new trigger component was added, the topological trigger (L1Topo). This component makes it possible to use more sophisticated requirements based on the topology and geometry of the event. The decision time for a Level-1 accept is $2.5\ \mu\text{s}$ [45].

The HLT is a software algorithm that runs on a large CPU farm of $\sim 40,000$ processor cores providing a trigger decision within 300 ms [46]. The HLT uses full granularity detector information in either the L1 determined RoIs or even the full event in some cases. The HLT selection algorithms are used not only for physics measurements but also for other purposes such as efficiency and performance measurements, detector monitoring and calibration.

Another challenge presented to the detector by the high luminosity conditions is *pile-up*. When multiple interactions occur in the same bunch crossing it is referred to as *in-time* pile-up, the primary measure for the level of pile-up is μ , the average number of inelastic interactions per bunch crossing. During Run II μ was typically around 30, but reached as high as 60 in certain periods, see Figure 3.2. Thus it is crucial for the detector systems to have sufficient granularity to distinguish the signals left by the particles produced from the overlapping pile-up collisions. This quantity μ is measured directly in the data via dedicated luminosity measurements provided primarily by the LUCID-2 [47] Cherenkov detector. LUCID consists of 16 photomultiplier tubes (PMTs) on each side of the ATLAS detector. Counts from these PMTs are integrated over 60 s periods called luminosity blocks to determine

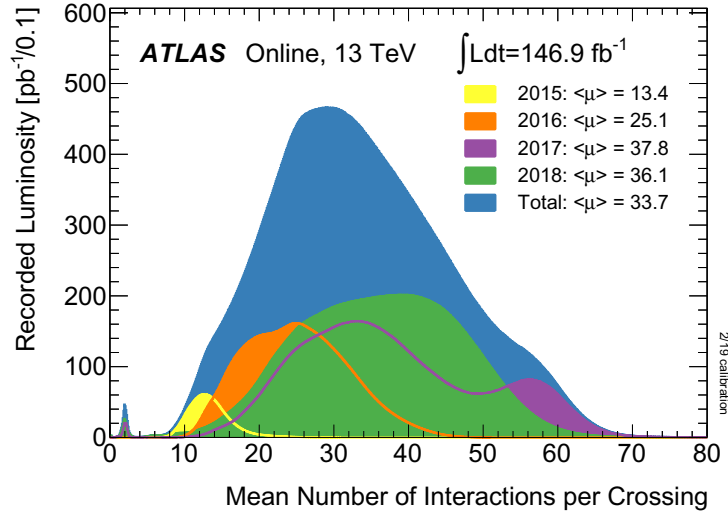


FIGURE 3.2: Mean number of interactions per bunch crossing in 13 TeV data collected by ATLAS from 2015 - 2018 [49].

the instantaneous luminosity. μ is determined from the instantaneous luminosity measurements via a calibrated conversion factor that accounts for the inelastic cross section (~ 80 mb) and LHC revolution frequency (11 246 Hz for pp collisions) [48].

A diagram depicting the layout of the ATLAS detector is shown in Figure 3.3 with each detector subsystem labelled and another depicting how different particle species interact with the detector subsystems is shown in Figure 3.4. The following three sections will describe each of the major detector systems, the inner tracking detector, the calorimeter and the muon spectrometer. Detector systems are typically split into *barrel* and *endcap* categories, with the barrel consisting of the central cylindrical barrel of the detector and the endcaps being the two forward regions nearer the beam pipe. The ATLAS detector is designed to have full coverage in ϕ and for $|\eta| < 5$ with precise coverage in the central (barrel) region of $|\eta| < 2.5$.

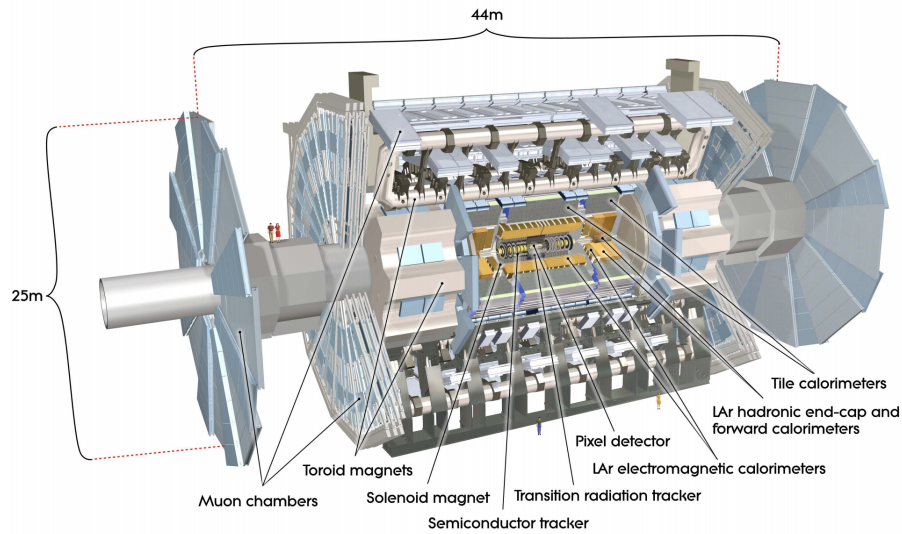


FIGURE 3.3: Diagram highlighting the detector systems of the ATLAS experiment [1].

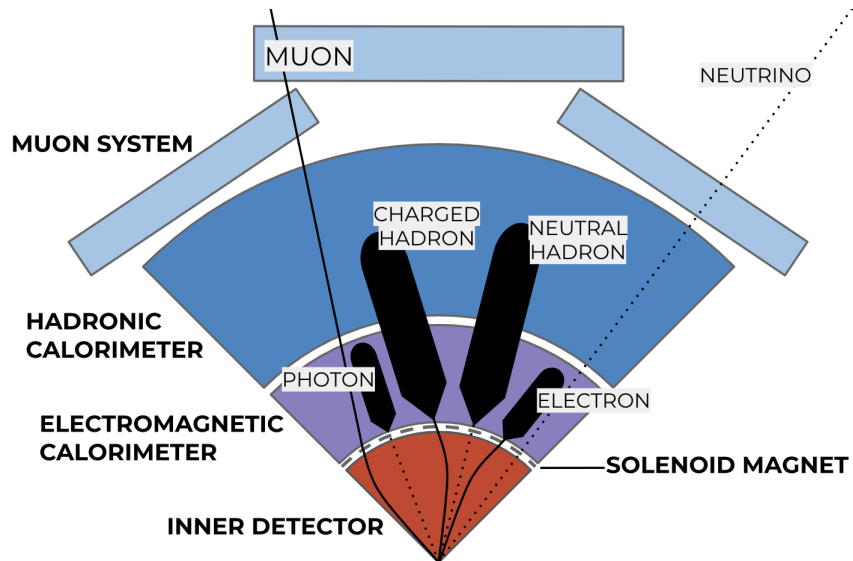


FIGURE 3.4: Diagram depicting the interactions of different species of particles with the ATLAS detector layers. Charged particles leave curved tracks in the inner detector, photons and electrons produce showers in the electromagnetic calorimeter, charged and neutral hadrons leave showers in the hadronic calorimeter. Muons are tracked by both the inner detector and muon system and neutrinos pass through undetected.

3.2.1 Inner detector

The inner detector (ID) is the tracking system used to detect particle trajectories and momentum in the region closest to the IP. The ID covers $|\eta| < 2.5$ and extends approximately 1 m in radius from the IP. It consists, in order of increasing distance from the IP, of silicon pixel detectors, silicon microstrip detectors (SCT) and the straw tube transition radiation tracker (TRT). The granularity of the detectors is highest in the layers closest to the IP and decreases as the tracks become spread over a larger area moving away from the collision point. Charged particles passing through the detector layers will deposit some energy that is detected and referred to as *hits*, tracking algorithms combine hits to determine particle trajectories.

The point at which an inelastic collision occurs is called a *primary vertex* (PV), in the high μ bunch crossings there will be many primary vertices. The precision of the ID is important for associating particle tracks to a primary vertex and resolving individual vertices, this is crucial for reducing the impact of pile-up. The main performance goal of the ID is a transverse momentum resolution of

$$\frac{\sigma_{p_T}}{p_T} = \left(0.05 \frac{p_T}{\text{GeV}} \oplus 1 \right) \%, \quad (3.1)$$

and a high- p_T track transverse impact parameter resolution, d_0^* , of 10 μm .

The barrel section of the ID depicted in Figure 3.5 extends to ± 80 cm in z consisting of concentric cylindrical detector layers. The endcap regions extend to $\sim \pm 2.8$ m and consist of detectors mounted on disks placed perpendicular to the beam axis (see Figure 3.6). The ID is surrounded by a relatively thin solenoid superconducting

* d_0 and z_0 are the transverse and longitudinal displacement of the track at the point closest to the primary vertex, respectively.

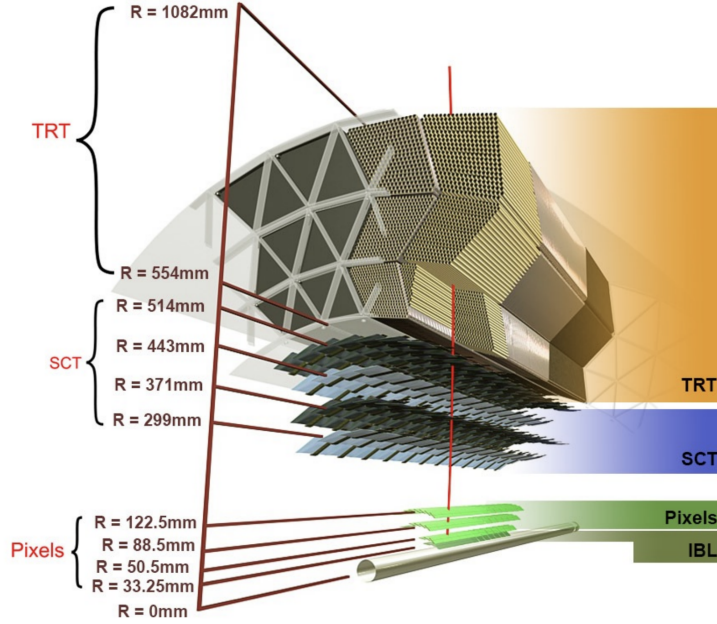


FIGURE 3.5: Layers of the ATLAS inner tracker barrel region including the insertable B-layer (IBL) [50].

magnet that produces a 2 T axial magnetic field to bend the trajectories of charged particles to determine their charge and transverse momentum.

The pixel [51] detectors originally consisted of three barrel layers, however in the shutdown proceeding Run II a forth layer was added closest to the beam pipe, the insertable B-layer (IBL) [52]. The IBL is designed to enhance the precision of the inner tracker particularly in vertex resolution and to mitigate the impact of radiation damage performance degradation. The layout of the pixel layers is such that four hits will be deposited by any charged particle within $|\eta| < 2.5$.

The three original pixel layers are comprised of 1744 identical pixel sensors with an external area of $19 \times 63 \text{ mm}^2$ made of $250 \mu\text{m}$ thick oxygenated n-type silicon wafers for radiation tolerance. The nominal pixel channel size is $50 \times 400 \mu\text{m}^2$ and there are ~ 80 million channels total across the pixel layers (67 million in the barrel

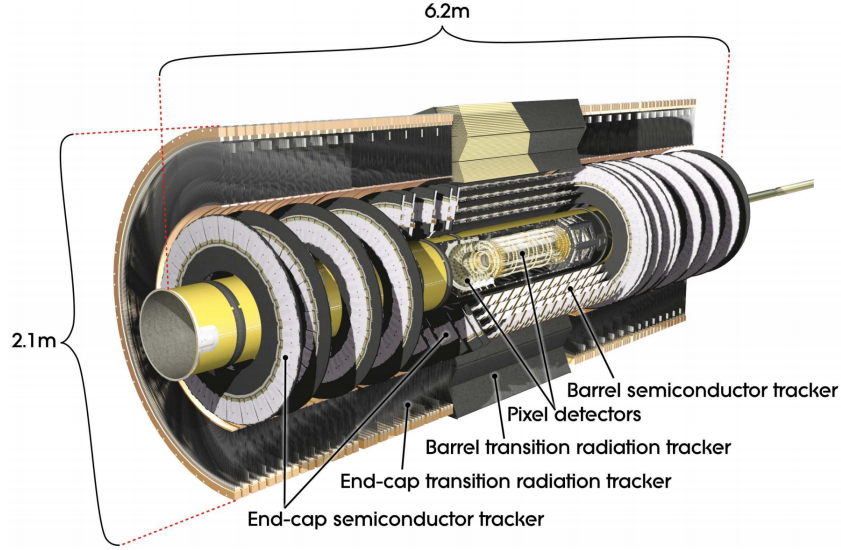


FIGURE 3.6: Diagram depicting the detectors of the ATLAS inner tracker [1].

and 13 million in the endcaps). The IBL is a fourth cylindrical barrel layer consisting of 224 pixel modules with individual pixel channels $50 \times 250 \mu\text{m}^2$ in size, the IBL added another 6 million readout channels.

The SCT [53] is comprised of 4088 modules each containing 768 silicon-strip channels constructed of $285 \mu\text{m}$ thick high-resistivity n-type bulk with p-type implants. These channels are very asymmetric rectangles of $80 \mu\text{m} \times 6 \text{ cm}$, thus providing a precise position measurement in only one direction. This limitation is overcome by pairing the strip modules back-to-back with a stereo angle of 40 mrad . In the barrel the strips are positioned approximately parallel to the beam axis and in the endcap disks they are positioned radially, both are generally perpendicular to the bending direction of the solenoid field. Four hits are expected in the SCT from particles originating from the IP.

The TRT [54] is comprised of thin-walled proportional drift tubes called *straws*.

A straw consists of a 4 mm diameter hollow aluminium tube with a kapton covered inner wall serving as the cathode held at -1530 V with respect to a $31\text{ }\mu\text{m}$ tungsten anode wire at the tube centre. During Run II the tubes were filled with a gas mixture of 70% Ar, 27% CO_2 and 3% O_2 , passing charged particles ionize the gas which drifts to the anode. The TRT barrel contains 52544 straws 144 cm in length positioned parallel to the beam axis, the two endcaps each contain 122880 straws 37 cm in length positioned radial to the beam axis. The straw technology was chosen because of the ability to instrument a relatively large area of the inner tracker at a reduced cost compared to a fully silicon based tracker. The resolution of a single TRT hit is worse than a silicon hit, however particles produced within $|\eta| < 2.0$ will cross 35-40 straws as they pass through the TRT which provides adequate tracking.

3.2.2 Calorimeter

The ATLAS calorimeter is composed of an inner electromagnetic calorimeter (ECAL) surrounded by a hadronic calorimeter (HCAL). Particle *showers* develop as the particles pass through the layers of the calorimeters. These layers alternate between *absorber* material which initiate the showers and *active* detector material to measure the shower energy. This design is referred to as a *sampling* calorimeter, the measured energy deposited in the active layers is calibrated to infer the shower energy accounting for the undetected energy deposited in the absorber layers. The ECAL measures the energy of electrons and photons deposited in EM showers and the HCAL measures hadronic showers.

As depicted in Figure 3.7, the ATLAS calorimeters fall into two technological categories: liquid argon (LAr) calorimeters with its namesake active material and

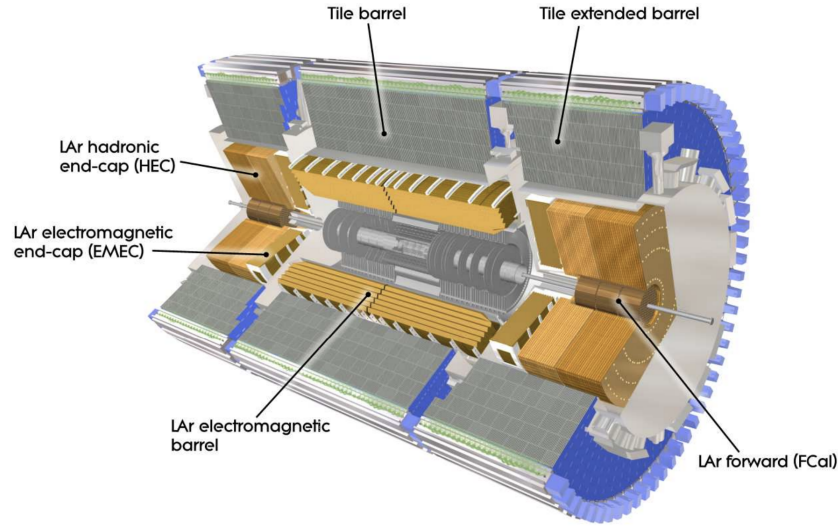


FIGURE 3.7: Diagram depicting the detectors of the ATLAS calorimeter [1].

various passive absorber materials and, tile calorimeter (TileCal) using steel absorption layers and scintillating tiles as the active medium. The ECAL is exclusively LAr calorimeters and the HCAL consists of LAr calorimeters in the high radiation forward areas and tile calorimeters elsewhere. LAr calorimeters are used in the regions exposed to high radiation due to their resilience. The calorimeters cover the range $|\eta| < 4.9$, with ECAL thickness of >22 radiation lengths ($X_0^\dagger$) and a HCAL thickness of ~ 10 interaction lengths ($\lambda^\ddagger$). This is sufficient for the calorimeter to effectively contain EM and hadronic particle showers and mitigate *punch-through* of charged particles into the muon system. Hadronic showers start in the electromagnetic calorimeter, which contains the first 2-3 interaction lengths.

[†]Radiation length, X_0 [$\text{g}\cdot\text{cm}^{-2}$], is a material characteristic, for each radiation length an electron will lose all but $1/e$ of its energy. The mean free path for pair production from a photon is also $\sim X_0$, thus the radiation length is the characteristic distance for the structure of an EM shower.

[‡]Interaction length, λ [$\text{g}\cdot\text{cm}^{-2}$], is also a material property characterizing the mean distance a hadronic particle travels before interacting. In general the interaction length is larger than the radiation length, thus hadronic showers are larger and more diffuse than EM showers.

The EM calorimeter consists of an EM barrel [55] covering $|\eta| < 1.475$, the EM endcap (EMEC) [55] covering $1.375 < |\eta| < 3.2$ and the forward calorimeter (FCal) [55] covering $3.1 < |\eta| < 4.9$. The EM barrel is 6.4 m long with an inner and outer radius of 2.8 m and 4 m, respectively. The EMEC is mechanically divided into an outer ($1.375 < |\eta| < 2.5$) and inner wheel ($2.5 < |\eta| < 3.2$), which together form a 630 mm thick wheel of inner and outer radius of 330 mm and 2098 mm, respectively. Both the barrel and endcap regions use lead as the absorber material, the lead plates are fashioned with an accordion geometry that naturally provides full ϕ coverage without any cracks. The EM calorimeter is divided into three layers that each serve different functions. Before the first layer is the presampler covering $|\eta| < 1.8$, which is designed to account for energy lost upstream of the calorimeter. Layer 1, closest to the IP, is highly segmented to determine precisely the position of photons that do not leave tracks in the inner detector. Layer 1 is $\sim 4 X_0$ thick and does not contain the shower, this is the function of the less granular but $\sim 16 X_0$ thick layer 2. Layer 3 is relatively thin and coarse, it is used to account for energy leaking beyond the ECAL. The EMEC inner wheel consist of only the first 2 layers. The ECAL contains a total of ~ 164 thousand readout channels (102k barrel, 62k endcap) and another ~ 11 thousand presampler channels.

The hadronic calorimeter consists of the TileCal [56] barrel covering $|\eta| < 1.7$ and the LAr hadronic endcap (HEC) [55] covering $1.5 < |\eta| < 3.2$. The barrel is mechanically divided into the 5.8 m central barrel covering $|\eta| < 1.0$ and two 2.6 m extended barrels covering $0.8 < |\eta| < 1.7$, each with an inner and outer radius of 2.28 m and 4.25 m, respectively. The scintillating tiles of the hadronic barrel produce UV light as they are traversed by charged particles, these are collected by photo

multiplier tubes located on two sides of the tile via wavelength shifting fibres. The TileCal is segmented into 64 modules each subtending 5.625 degrees in ϕ with ~ 10 thousand readout channels.

The HEC consist of four layers located directly behind the EMEC, unlike the EM LAr calorimeters the HEC absorbtion layers are flat and constructed from copper. Each HEC is divided into two cylindrical wheels (HEC1 and HEC2) of outer radius 2030 mm with a combined thickness of 1818 mm. The inner radius is 372 mm for a small segment of the HEC nearest the IP and 475 mm otherwise. HEC1 is closer to the IP and features finer segmentation compared to HEC2, together containing 5632 readout channels.

Each FCal consists of three 45 cm deep (z-direction) cylindrical LAr modules of 47 cm outer radius. The front face of the FCal is recessed by 1.2 m with respect to the EMEC and has a much shorter physical depth, therefore the FCal is designed to be highly dense to contain particle showers. FCal1 is designed to contain EM showers and uses copper as the absorber, FCal2 and FCal3 use mainly tungsten as the absorber to contain the lateral spread of hadronic showers. The FCal uses a unique detector geometry consisting of cylindrical electrodes positioned parallel to the beam axis and supported by copper endplates. An electrode consist of a metal rod (copper or tungsten) and coaxial copper tube separated by a radiation-hard plastic fibre wound around the rod, the gap formed by the fibre is filled with active material (liquid argon). The electrodes are grouped into 3524 readout channels.

The material depth in terms of interaction lengths (λ) for the entire calorimeter is shown in Figure 3.8. The calorimeter provides $\sim 10\lambda$ across the full η range effectively containing hadronic showers, important for jet energy resolution.

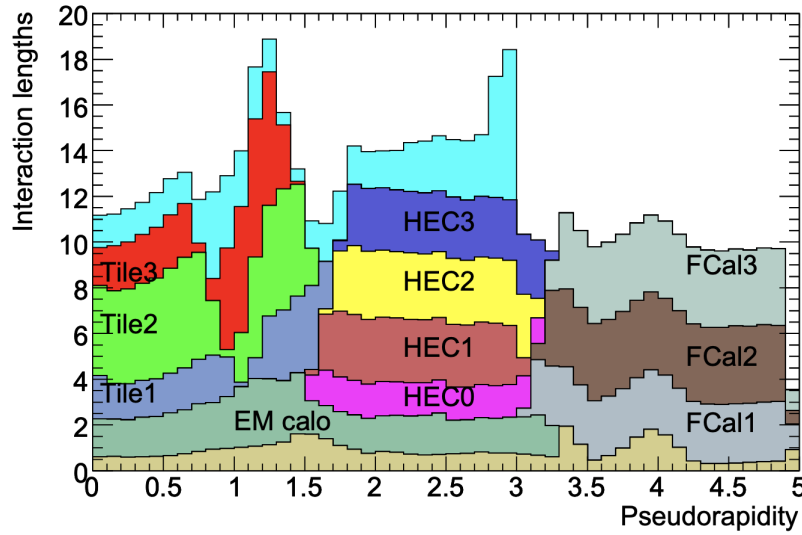


FIGURE 3.8: Cumulative material depth, in interaction lengths as a function of $|\eta|$. The bottom segment represents the material in front of the ECAL and the top segment accounts for the total material present before the muon spectrometer (for $|\eta| < 3.0$) [1].

3.2.3 Muon spectrometer

The muon spectrometer (MS) [57] is the largest and outermost detector system, it is located outside the calorimeter in the region where only muons and neutrinos are expected to reach. The MS consists of chambers placed in and around the toroid magnets in the barrel and of radial disks of chambers in the endcaps as depicted in Figure 3.9. Opposite to the solenoid field present in the inner tracker, the toroid field is designed to bend muons in the η direction, making η the precision coordinate. The toroid magnet system consists of: a main set of eight large air-core superconducting magnets arranged symmetrically outside the calorimeter barrel producing a ~ 0.5 T field and, two smaller sets of eight similar magnets in the endcaps producing a ~ 1.0 T field (see Figure 3.10). Particle tracking chambers cover $|\eta| < 2.7$ and the MS is also designed to trigger on particles in the range $|\eta| < 2.4$.

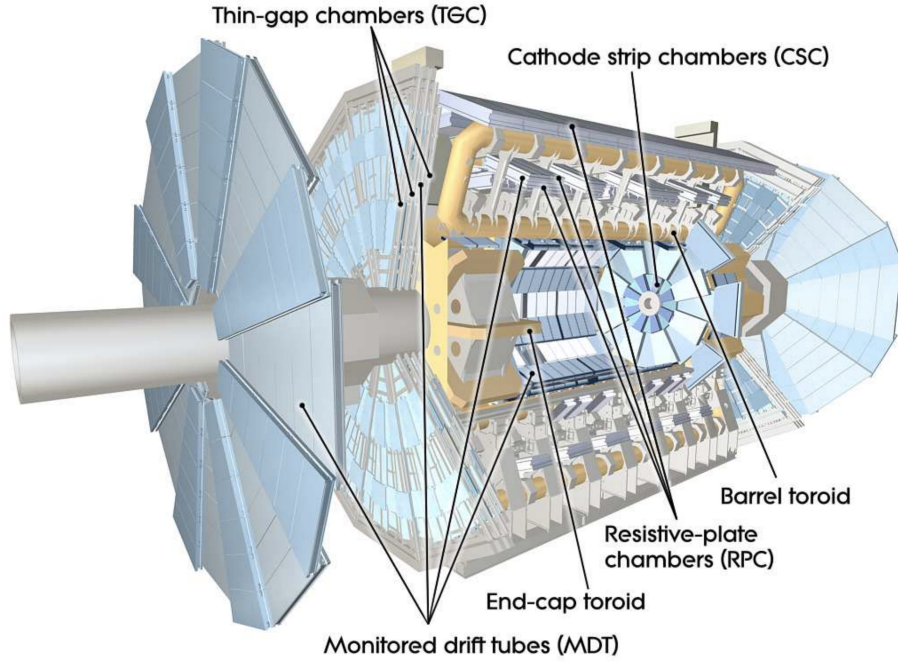


FIGURE 3.9: Diagram depicting the detectors of the ATLAS muon spectrometer [1].

The design goal of the MS is to determine muon charge with high accuracy and to measure the p_T of a 1 TeV muon with 10% resolution. The spectrometer performs well over a wide range of muon momentum from ~ 3 GeV to ~ 3 TeV, with momentum resolution gradually degrading as momentum increases due to track curvature decreasing.

The MS chambers consists of four detector technologies: monitored drift tubes (MDT) and cathode strip chambers (CSC) responsible for precision tracking and, resistive plate chambers (RPC) and thin gap chambers (TGC) responsible primarily for triggering and to a lesser extent tracking. The chambers are arranged in eight octants to reflect the symmetry of the toroid magnets, with each octant subdivided into a large and small sector that overlap slightly in ϕ . The barrel chambers are

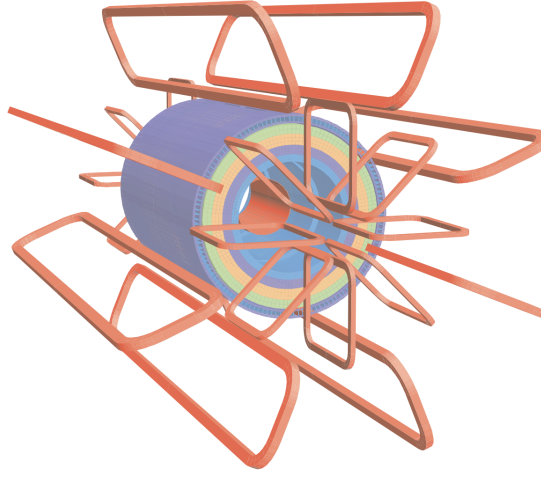


FIGURE 3.10: Geometry of the toroid magnet coils, the endcap toroids are rotated 22.5° relative to the barrel toroids. The central cylinder is a model of the tile calorimeter [1].

arranged in three concentric cylindrical shells at radii of approximately 5 m, 7.5 m and 10 m. There is a gap in chamber coverage in the centre of the detector ($|\eta| \approx 0$) to allow for services to the inner systems of the detector. The size of the gap varies but is at most 1-2 m. The endcap chamber wheels are located at distances of $|z| \approx 7.4$ m, 10.8 m, 14 m and 21.5 m from the IP, the position of the various MS detectors is shown in Figure 3.11.

The vast majority of tracking is handled by MDT chambers, chosen for the combination of high measurement accuracy and simplicity of construction. MDT chambers consist of three to eight layers of 30 mm diameter tubes filled with a mixture of 93% Ar and 7% CO₂, while containing a central tungsten anode wire. The tubes are aligned parallel to the beam providing a $35 \mu\text{m}$ resolution in η (per chamber) and a very coarse ϕ position measurement as the tubes run the full length of the chambers. The Run 1 configuration of MDT chambers consisted of 1150 chambers with 354000 readout channels. For Run II, 14 additional half sized MDT

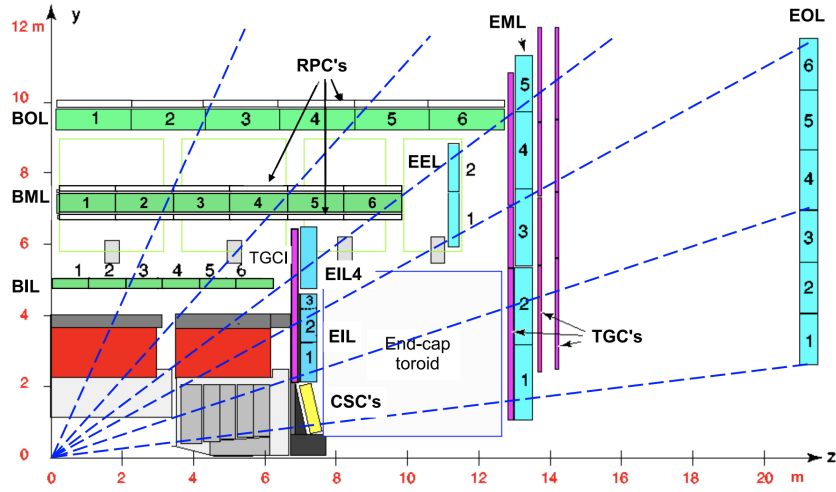


FIGURE 3.11: Positions of the muon spectrometer detectors in the yz plane, RPCs, TGCs and CSC are labelled by name, the blue and green chambers labelled with numerical values are all MDTs [1].

chambers were added to fill gaps in acceptance in certain areas related to detector support structures and services, these chambers contain ~ 6400 additional readout channels [58].

In the endcaps, the innermost layer of tracking chambers covering $2.0 < |\eta| < 2.7$ are CSCs due to their higher rate capability necessary in the forward region. Thus, there are only 32 CSCs with a total of 31000 readout channels. These are multiwire proportional chambers with two cathode readout planes, segmented into strips in orthogonal directions to measure both η and ϕ coordinates. Each CSC consists of four 5 mm gaps producing four independent η , ϕ measurements, filled with a mixture of 80% Ar/20% CO₂ and containing a plane of 30 μm tungsten wires. The strip pitch and width varies in the orthogonal directions leading to a position resolution per chamber of approximately 40 μm in η and 5 mm in ϕ .

The barrel trigger chambers for $|\eta| < 1.05$ are RPCs and TGCs for the endcaps

covering $1.5 < |\eta| < 2.4$. Both RPCs and TGCs can deliver a time signal with a resolution of 15-25 ns, providing the capability to distinguish the bunch crossing a given muon track originates from with $\geq 99\%$ efficiency. The trigger chambers also provide a relatively coarse position measurement with a resolution of 2-7 mm (TGC) and 10 mm (RPC) in both η (R/z) and ϕ directions. The position measurement provides a coarse momentum resolution for use in triggering, and the ϕ measurements of the trigger chambers are much more precise than that of the MDTs.

In the barrel the RPCs are positioned in three concentric cylindrical shells attached to the outside surfaces of the MDT chambers, with equal dimension and a common support structure. The RPC is a gaseous parallel electrode-plate detector consisting of two 2 mm gas gaps filled with a gas mixture of $\text{C}_2\text{H}_2\text{F}_4/\text{Iso-C}_4\text{H}_{10}/\text{SF}_6$ (94.7/5/0.3) between two phenolic-melaminic plastic laminate resistive plates. The plates are read out by capacitive coupling to metallic strips with 23-35 mm pitch mounted on the outside of the plates, positioned in the z direction on one side and ϕ on the other. There are 606 RPCs with 373000 readout channels.

The TGCs are positioned in four layers around the endcaps, one layer preceding the small wheel of CSCs and MDT chambers at $|z| \approx 7.4$ m, one layer preceding the large MDT wheel at $|z| \approx 14$ m and two layers after this wheel. The TGCs are multiwire proportional chambers like the CSC, but optimized fast readout over position resolution, consisting of 2.8 mm thin gas gaps filled with a 55% CO_2 /45% n-pentane mixture and 50 μm tungsten wires. There are 3588 TGCs with 318000 readout channels.

Chapter 4

Event reconstruction and simulation

The reconstruction of collisions recorded with the ATLAS detector is a crucial step for any physics measurement. Also important is the simulation of the collisions and of the detector response; experimental physicist devote a lot of time to these aspects of the experiment.

Personally, I have contributed to the simulation of the *new muon small wheel*, an upgrade to the muon system set to be installed during the 2019-2021 long shutdown of the LHC. My contribution focused on the simulation of the digitization of detector signals responding to muon energy deposits and involved validation of the simulations against test beam data collected at Fermilab. I have also contributed to the jet/missing transverse energy combined performance group. I studied the impact on jet energy scale and jet mass response of a variety of jet reconstruction algorithms and so-called *jet grooming* techniques [59]. These grooming techniques are mainly designed to reduce the pileup dependence of jet observables and to improve jet calibration performance.

This chapter outlines the details of the various MC simulations that were analysed. The details of the object selection are also described as well as corrections

applied to the simulation to provide results that better represent the observed data.

4.1 Data and Monte Carlo samples

4.1.1 Data

The LHC and ATLAS detector generally performed very well during Run II in terms of data collection rate and data quality. The total integrated luminosity was ~ 6 times what was collected in Run I. As is shown in Figure 4.1, each year of data taking between 2015–2018 outperformed the previous year in terms of total integrated luminosity delivered by the LHC and recorded by ATLAS. This volume of data makes new searches possible and provides unprecedented statistical precision for full Run II measurements.

This analysis uses the full dataset collected during Run II of the LHC, corresponding to a total integrated luminosity of $139.0 \text{ fb}^{-1} \pm 1.7\%$ [60]. The mean number of interactions per bunch crossing, μ , differs significantly between each year of data taking as shown in Figure 3.2. The integrated luminosity and μ of each year of data taking is summarized in Table 4.1. The data collected in 2016 suffers from a small excess in the number of high energy calorimeter clusters in the region around $|\eta| \sim 2.9$ [61]. The affected events are flagged as bad and can thus be vetoed as part of the event selection. Taking this into account, the total luminosity of the data used is $138.4 \text{ fb}^{-1} \pm 1.7\%$.

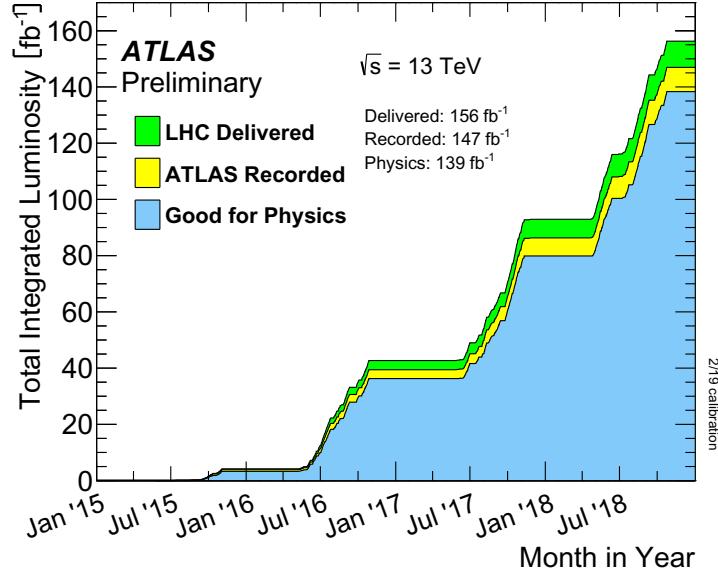


FIGURE 4.1: Total integrated luminosity versus time delivered to ATLAS by the LHC (green), recorded by ATLAS (yellow) and certified as good quality data (blue) from pp collisions during Run II [49].

Year	Luminosity [fb^{-1}]	$\langle\mu\rangle$
2015	3.2	13.4
2016	33.0	25.1
2017	44.3	37.8
2018	58.5	36.1
Total	139.0	33.7

TABLE 4.1: Summary of the integrated luminosity collected during each year of data taking by the ATLAS detector during Run II of the LHC [49].

4.1.2 MC simulation

MC generators are used to simulate events produced by the EW- Zjj signal process and background processes (dominated by strong- Zjj). The resulting event samples capture the characteristics of the different physical processes, and are used in the data analysis to estimate the background and optimize the selection criteria. MC samples are used to evaluate systematic uncertainties by introducing variations into the samples and observing their effect. MC simulation is also used in the procedure to extract the EW- Zjj signal from the data to produce the detector-level results. Finally, the samples are used in the procedure that corrects for detector inefficiency and resolution (Section 6.7) to produce the final particle-level results.

To better represent the varying pileup conditions of the different years of data taking (see Figure 3.2), the MC simulation were generated in several *MC campaigns*. The *MC16a* campaign matches the data collected in 2015 and 2016. The 2015 and 2016 data taking periods are combined as the 2015 dataset is too small to warrant its own separate campaign. *MC16d* corresponds to 2017 data and *MC16e* corresponds to 2018 data. These three campaigns, and the corresponding data, are treated separately in this analysis. Physics *objects* (jets, leptons) are reconstructed and calibrated according to the recommendations of the ATLAS combined performance groups which often differ year-to-year.

Accurate modelling of the EW- Zjj and strong- Zjj processes is very challenging and crucial for this analysis. As such, three different MC simulations that use different simulation techniques have been analyzed for each process. The other background processes, such as top and diboson, are of lesser impact. Only a single MC sample is considered for each of these processes. In what follows here is the

details of the MC simulations, these are summarized in Table 4.2.

The default simulation of EW- Zjj was produced with POWHEG-BOX v1 [62–64] using the CT10nlo [65] PDFs, and is accurate to NLO in perturbative QCD. This sample was produced with the “VBF-approximation” [66]. The parton shower, hadronization and underlying event were simulated with PYTHIA 8.186 using the AZNLO [67] set of tuned parameters. The EVTGEN program [68] was used for the properties of the bottom and charm hadron decays. This sample is referred to as POWHEG+PY8 EW- Zjj production.

The second EW- Zjj sample was produced with HERWIG 7.1.5 [69, 70], also using the VBF approximation and accurate to NLO in perturbative QCD, with loop-amplitudes provided by VBFNLO 3.0.0 [71]. The MMHT2014LO PDF set [21] was used and the parton shower, hadronization and underlying event was also simulated with HERWIG 7.1.5 using the default set of tuned parameters. EVTGEN was used for the properties of the bottom and charm hadron decays. This sample is referred to as HERWIG+VBFNLO EW- Zjj production.

The third EW- Zjj sample was produced with SHERPA 2.2.1 [72], again in the VBF approximation with LO accuracy in the matrix element calculation with up to two additional parton emissions. The NNPDF3.0nnlo PDFs [73] were used and the MEPS@LO prescription [24] was used to merge the matrix element with the SHERPA parton shower. Hadronization and the underlying event were simulated with SHERPA using the default set of tuned parameters. This sample is referred to as SHERPA EW- Zjj production. The SHERPA prediction significantly underestimates the measured differential cross-sections, due to a non-optimal setting of the colour flow [74].

The first strong- Zjj sample was produced with SHERPA 2.2.1, using the Comix [75] and OpenLoops [76, 77] libraries to generate predictions of $Z + n$ partons ($n \leq 4$). The predictions are accurate to NLO for up to two partons in the final state and to LO for three or four partons. The different matrix element level calculations were merged into an inclusive sample and matched to the parton shower with an improved CKKW procedure [24, 25], extended to NLO accuracy with the MEPS@NLO prescription [78]. The NNPDF3.0nnlo PDF set was used and the prediction was normalized to a NNLO prediction for inclusive Z boson production [79]. Hadronization and the underlying event were also simulated with SHERPA using the default set of tuned parameters. This sample is referred to as SHERPA strong- Zjj production.

The second strong- Zjj sample was produced with MADGRAPH5_aMC@NLO [65], and is also accurate to NLO for up to two partons in the final state and LO for three or four partons. The NNPDF2.3nlo PDF set [80] was used and the matrix element and parton shower were merged with the FxFx prescription [28]. The MADGRAPH generator was interfaced to PYTHIA 8.186 to provide the parton shower, hadronization and underlying event using the A14 set of tuned parameters [81]. The sample is also normalized to the same NNLO prediction as the SHERPA sample. EVTGEN was used for the properties of the bottom and charm hadron decays. This sample is referred to as MG5_NLO+PY8 strong- Zjj production.

The third strong- Zjj sample was also produced with MADGRAPH5_aMC@NLO, however for this sample the accuracy is LO up to four partons in the final state. The NNPDF3.0lo PDF set was used and the matrix element and parton shower were merged with the CKKW-L procedure [26, 27]. Again, PYTHIA 8.186 provides the parton shower, hadronization and underlying event using the A14 set of tuned

parameters and the sample is normalized to the same NNLO prediction. This sample is referred to as MG5+PY8 strong- Zjj production.

Diboson final states are produced with SHERPA at NLO accuracy for up to one parton in the final state, and at LO for two or three partons. These samples are subdivided by the decay modes of the diboson pairs, two such modes (ZV with $V \rightarrow jj$ and $Z \rightarrow \ell\ell$ where $V = W, Z$) are part of inclusive Zjj production, the rest are not. The NNPDF2.0nnlo PDF was used and the virtual corrections were taken from OpenLoops. The parton shower merging was performed with the MEPS@NLO algorithm and the hadronization and underlying event were simulated with SHERPA using the default set of tuned parameters. These two samples are referred to as diboson Zjj .

Non- Zjj samples are included in the analysis to study the contribution of background processes to the $\ell^+\ell^-jj$ final state. In these processes the opposite sign leptons do not originate from the decay of a Z boson, but by chance the invariant mass of the dilepton systems falls near the Z peak. The first category of these backgrounds is diboson production other than the $ZV \rightarrow \ell^+\ell^-jj$ processes, these are produced with SHERPA as described above for the diboson Zjj samples.

The second category are events containing a single top quark or a top-antitop ($t\bar{t}$) pair, these are simulated with POWHEG-BOX v2 at NLO accuracy using the hqv program [82]. The NNPDF3.0nnlo PDF set was used and the generator was interfaced to PYTHIA 8.230 to provide the parton shower, hadronization and underlying event using the A14 set of tuned parameters.

The third category are events from W +jets and $Z \rightarrow \tau\tau$ processes. These were included for completeness but the contribution from these processes to the VBF

Process	Generator	Matrix element accuracy	PDF set	Shower and hadronization	Parameter set
EW- Zjj	POWHEG-BOX v1	NLO	CT10nlo	PYTHIA8+EVTGEN	AZNLO
	HERWIG7 + VBFNLO	NLO	MMHT2014lo	HERWIG7+EVTGEN	default
	SHERPA 2.2.1	LO(2-4j)	NNPDF2.0nnlo	SHERPA	default
strong- Zjj	SHERPA 2.2.1	NLO(0-2j) LO(3-4j)	NNPDF3.0nnlo	SHERPA	default
	MADGRAPH5_aMC@NLO	NLO(0-2j) LO(3-4j)	NNPDF2.3nlo	PYTHIA+EVTGEN	A14
	MADGRAPH5	LO(0-4j)	NNPDF3.0lo	PYTHIA+EVTGEN	A14
VV	SHERPA	NLO(0-1j) LO(2-3j)	NNPDF3.0nnlo	SHERPA	default
$t\bar{t}$, single top	POWHEG-BOX v2 hvq	NLO	NNPDF3.0nnlo	PYTHIA+EVTGEN	A14
W +jets, $Z \rightarrow \tau\tau$	POWHEG-BOX v1	NLO	CT10nlo	PYTHIA8+EVTGEN	AZNLO

TABLE 4.2: Summary of generators used for simulation. The details and the corresponding references are provided in the body of the text. In the final column, ‘default’ refers to the default set of tuned parameters provided with the event generator.

topology is negligible. Samples are produced with POWHEG-BOX v1 at NLO accuracy. The parton shower, hadronization and underlying event were simulated with PYTHIA 8.186 using the AZNLO set of tuned parameters. EVTGEN was used for the properties of the bottom and charm hadron decays.

The simulated events were passed through the GEANT4 [83] simulation of the ATLAS detector [84] and reconstructed using the same procedure as used for the data. Pileup is simulated by overlaying inelastic pp interactions generated by PYTHIA8 [85], with the A3 tune [86] and NNPDF2.3LO PDF set. These inelastic interactions are added to the simulated samples of each MC campaign such that the μ distribution of the simulation matches the corresponding data period. The samples containing physics objects reconstructed in the detector are said to be at *detector-level*.

Two samples, HERWIG+VBFNLO EW- Zjj and MG5_NLO+PY8 strong- Zjj exist only at particle-level. Approximate detector-level predictions for these samples

are obtained by tuning the MC event weights (reweighting) of another MC sample (one that does have a detector-level prediction) to match kinematic distributions of the particle-level only prediction. Samples where such a particle-level reweighting has been applied are labelled with a prime symbol. The detector-level prediction of HERWIG+VBFNLO EW- Zjj is obtained by reweighting POWHEG+PY8 EW- Zjj to match HERWIG+VBFNLO at particle-level, this sample is referred to as HERWIG+VBFNLO'. Similarly, MG5_NLO+PY8' is obtained by reweighting MG5+PY8 strong- Zjj at particle-level. The procedure used to reweight the samples is described in Appendix A.

4.2 Event and trigger requirements

The ATLAS data is susceptible to various detector related issues. The following event level requirements are placed on the Run II data to ensure all detectors were fully operational:

- Be part of the standard ATLAS good run list*,
- Do not suffer from problems with tile calorimeter, liquid argon system or silicon tracker,
- Do not suffer from data acquisition problems, specifically missing information due to a restart of the timing, trigger and control system.

These detector-level problems affect a small number of events, the impact on the total integrated luminosity is negligible.

*The good run list (GRL) is a file that specifies whether the data taking conditions during each 1-2 minute interval (*luminosity block*) are acceptable for physics analysis.

Events are collected by either a two muon or two electron trigger (a dilepton trigger). The centrally recommended triggers are used. These differ for each data period and they are listed in Table 4.3. All triggers used are un-prescaled, meaning that all events that fulfil the trigger requirements were recorded. The trigger p_T thresholds were increased over the course of Run II. This was to counteract the increased rate of pileup and trigger rates as the instantaneous luminosity was ramped up.

The first part of the trigger names in Table 4.3, HLT, denotes that the required triggers are all high level triggers. The second part of the names indicate the p_T requirement in GeV. For example, 2e12 means two electrons with $p_T > 12$ GeV. The last piece specifies the selection requirements applied to the electrons or muons. The electron triggers are likelihood based triggers. This means that the selection is tuned such that the physics objects selected by the trigger have a certain likelihood of the observed signal being due to a real electron (signal) versus due to a hadron (background). This likelihood tuning presents a trade-off between the *efficiency* of selecting electrons and the *purity* of the selected objects. A *tight* likelihood requirement will be more pure, and a *loose* requirement will have higher efficiency at selecting the signal objects.

The muon triggers are similar in each period, with only the p_T requirement changing. The first mu## specifies the p_T requirement for a single L1 muon. The second muon does not need to be present at L1 (noL1). The efficiency is increased by scanning the full detector in the HLT requiring a second muon with $p_T > 8$ GeV (mu8).

Data year, MC period	Electron channel	Muon channel
2015	HLT_2e12_1hloose_L12EM10VH	HLT_mu18_mu8noL1
2016, mc16a	HLT_2e17_1hvloose_nod0	HLT_mu22_mu8noL1
2017 & 2018, mc16d & mc16e	HLT_2e24_1hvloose_nod0	HLT_mu22_mu8noL1

TABLE 4.3: Overview of the dilepton triggers. 1h specifies that these are likelihood based triggers with the loose or very loose (vloose) likelihood thresholds. nod0 indicates that no d_0 information is used when determining the likelihood. In 2015, the electron trigger specifies L12EM10VH. First, this means that two L1 ECAL trigger objects with $p_T > 10$ GeV are required, specified by L12EM10. Next, V indicates that the p_T requirement varies with η to account for energy loss. Finally, H denotes hadronic core isolation, meaning that the central core of the RoI of the HCAL for the trigger object must be less than or equal to some isolation threshold.

4.3 Object reconstruction and definition

In order to study fundamental interactions of the universe, one needs to contextualize the detector signals in terms of the underlying particles that created them. In a physics analysis, a physics *object* is a collection of detector signals interpreted as a final state particle, or as a jet of particles originating from a final state parton. The object will be identified as either a specific species of particle (typically electron, photon, muon or tau) or as a jet. The jet object can also contain flavour information about the species of parton it most likely originated from. Reconstruction algorithms combine information from the inner detector, calorimeters and muon spectrometer to identify the most likely underlying physics responsible for the observed signals. When reconstructing objects it is first important to know precisely the primary vertex they originate from.

4.3.1 Tracks and vertices

In the context of the ATLAS experiment, a particle track is a collection of energy deposits (*hits*) in the silicon and TRT layers of the inner detector that have been interpreted as the path travelled by a charged particle originating from a pp collision. Particle tracks in ATLAS are defined by their position and momentum at the point of closest approach to the beam spot. The longitudinal and transverse distances to the beam spot at closest approach are the track's longitudinal (z_0) and transverse (d_0) *impact parameters*. The position and momentum of the track can be extrapolated to any other point along its path based on knowledge of the detector material and magnetic field configuration.

Once found, the particle tracks are used to find the interaction vertices in the recorded events. The track and vertex finding algorithms use pattern recognition to make sense of the large number of inner detector hits present in a typical event with Run II pileup conditions. An extreme example of track and vertex reconstruction in a $Z \rightarrow \mu\mu$ candidate event with 65 additional pileup vertices is shown in Figure 4.2.

Tracks are reconstructed starting from track seeds built from three space point measurements in the pixel and/or SCT layers. Track seeds are extrapolated to all layers of the silicon inner tracker with a combinatorial Kalman Filter algorithm [88] to form track candidates. At this stage there are many so called *fake* track candidates and many hits are shared between candidates. A fake track candidate consists of detector hits that align by random coincidence due to the high density of the energy deposits. Candidate tracks are ranked with a score, which takes into consideration the number of associated hits, the type of hits (the more precise pixel hits are preferred over SCT hits) and the quality of the track fit. Shared hits are assigned

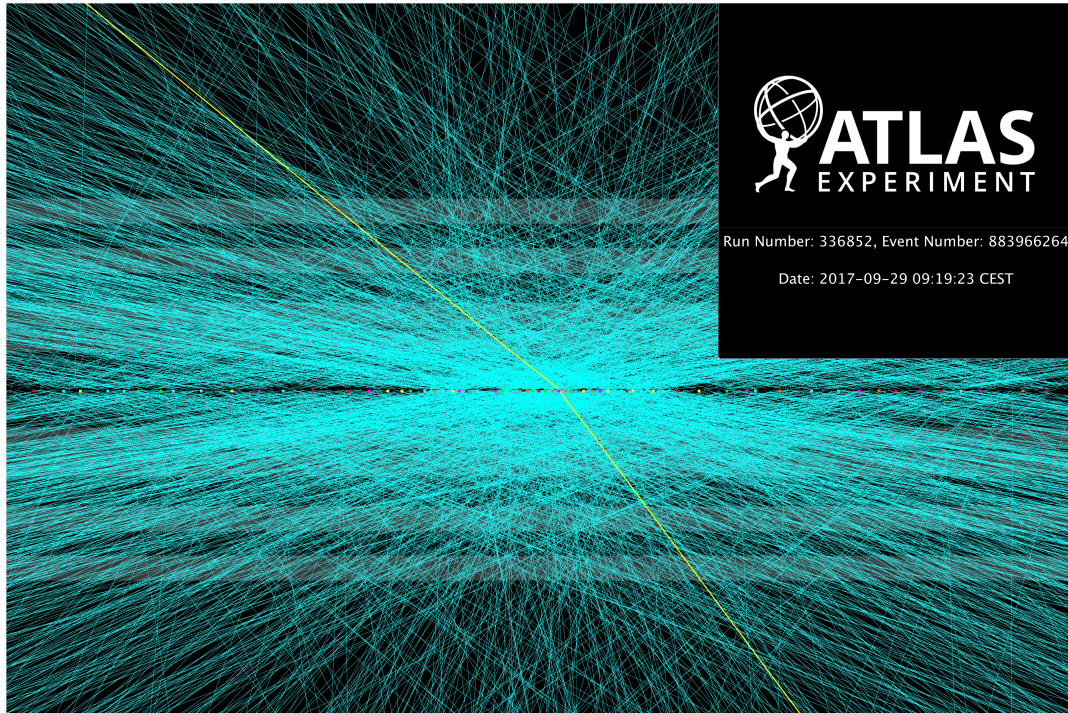


FIGURE 4.2: Reconstructed tracks and vertices from a $Z \rightarrow \mu\mu$ candidate event recorded in 2017. 65 pileup vertices are reconstructed in addition to the hard scatter vertex. All reconstructed tracks with $p_T > 100 \text{ MeV}$ are displayed and the two muon tracks are shown in yellow [87].

to the track with the highest score to resolve ambiguities and remaining tracks are dropped if they do not meet a track score threshold. Finally, the tracks are extended by adding matched hits in the TRT [89].

The identified tracks are then used for the vertex finding procedure. Of interest here are only the *primary* vertices at which a pp collision occurred. Events will also contain *secondary* or *displaced* vertices from the decay of short lived states, such as B mesons. Secondary vertices are identified with a separate dedicated procedure. For primary vertex finding, only tracks satisfying the “Tight” selection criteria [90] are used, this selection is optimized to reject fake tracks. The seed position for the vertex finding algorithm is taken from the xy position of the beam spot in the transverse plane. In the longitudinal direction the position is determined from the mode of the z_0 impact parameter of all selected tracks. With the seed determined, an iterative χ^2 minimisation procedure is used to find the optimal vertex position based on the position of closest approach of the tracks with tracks weighted based on their compatibility with the vertex position.

Once the optimal vertex position is found the tracks that are not compatible with the vertex position within seven standard deviations are removed from the vertex candidate and returned to the pool of unused tracks. The procedure is repeated until all tracks are associated to a vertex candidate or until no additional vertex can be found among the remaining tracks. Vertex candidates with at least two associated tracks are kept [91]. The vertex with the highest $\sum p_T^2$ is selected as the hard scatter vertex of the event, the remaining primary vertices are considered as pileup vertices.

4.3.2 Electrons

An electron is defined as an object consisting of a cluster built from energy deposits in the EM calorimeter and a matched track (or tracks) in the inner detector. In Run II the electron reconstruction algorithms used dynamic, variable-sized calorimeter clusters, called superclusters [92]. Electrons are identified using a likelihood-based method with different optional *working points*, which are a set of quality requirements applied depending on the needs of individual analyses. For electrons produced in a typical electroweak process, like Z boson production, the electrons identification efficiencies are, on average, 93%, 88% and 80% for the “Loose”, “Medium” and “Tight” working points, respectively [92]. Electrons selected for this analysis are required to satisfy the Medium likelihood working point, which trades reduced efficiency for improved background rejection relative to the Loose working point.

The electrons are also required to have $|\eta| < 2.47$, in order to fall within the coverage of the inner tracker. There is also a veto region between $1.37 < |\eta| < 1.52$ where electrons are removed as this regions corresponds to the calorimeter crack, a region with poor energy resolution due to the transition between the barrel and endcap calorimeters. Electrons are also required to have $p_T > 25 \text{ GeV}$ and satisfy the “Gradient” isolation working point. Isolation places requirements on the maximum energy deposited in a region around the reconstructed calorimeter cluster or inner detector track associated to the electron. The Gradient working point uses a cone of radius $R = 0.2$ around the object and optimizes selection criteria to target an efficiency rather than a fixed energy threshold, the isolation efficiency is 90% for an

electron with $p_T > 25$ GeV and 99% for $p_T > 60$ GeV [92]. Finally there are *track-to-vertex* association requirements on the transverse (d_0) and longitudinal (z_0) impact parameters:

- $|d_0/\sigma_{d_0}| < 5$,
- $|\Delta z_0 \sin \theta| < 0.5$ mm.

4.3.3 Muons

A muon is defined as an object consisting of an inner detector track matched to a muon spectrometer track. Track reconstruction is first performed independently in the ID and MS and the information is combined to form the muon objects used in analysis. Muon reconstruction is also performed according to a variety of muon quality working points, that are optimized to suppress background from pion and kaon decays [93]. The muons used in this analysis must satisfy the “Medium” quality requirements, which means that the muons must have ≥ 3 hits in at least two MDT layers[†] and a q/p significance[‡] of less than seven. The Medium quality requirement has an efficiency of 96.1% at selecting muons with $20 \text{ GeV} < p_T < 100 \text{ GeV}$ originating from W boson decays [93].

The ID tracks associated to the reconstructed muons must satisfy the following requirements:

[†]Muons with $|\eta| < 0.1$ are allowed to have hits on at least one MDT layer, but no more than one MDT hole layer.

[‡] q/p significance discriminates between prompt muons and background muons from hadron decay. It is defined as the absolute value of the difference between the ratio of the charge and momentum of the muons measured in the ID and MS divided by the sum in quadrature of the corresponding uncertainties [93].

- Number of pixel hits + number of crossed dead pixel sensors > 0 ,
- Number of SCT hits + number of crossed dead SCT sensors > 4 ,
- Number of pixel holes + number of SCT holes < 3 ,
- A successful extension into the TRT for $0.1 < |\eta| < 1.9$ [93],

where a *dead* sensor is known to be inactive and a hole is an active sensor traversed by the track but containing no hits. The hole must lie between two hits assigned to the track. Selected muons have $p_T > 25 \text{ GeV}$, $|\eta| < 2.4$ and must satisfy the same Gradient isolation working point as the electrons. Finally, similar to electrons, there are track-to-vertex association requirements:

- $|d_0/\sigma_{d_0}| < 3$,
- $|\Delta z_0 \sin \theta| < 0.5 \text{ mm}$.

4.3.4 Jets

As discussed in Section 2.3, jets are collimated groups of hadrons that are proxies for the final state quarks and gluons produced in collisions. For practical purposes, a *jet algorithm* is employed to reconstruct jets according to a well defined procedure. The procedure starts from a set of input particles that will become the *jet constituents*, these constituents can be calorimeter clusters or particle tracks. The algorithm systematically determines which constituents belong to which jet and how the constituent properties, such as transverse momentum, are combined to define the jet properties. The jet algorithm will depend on a set of parameters that are tuned

to suit the needs of the experiment. An important parameter is the *jet radius* (R) that characterizes the angular extent of the jets in the $y - \phi$ plane. Particles separated by a distance greater than R will not be part of the same jet.

The jet algorithm also specifies the *recombination scheme*, this is how the kinematic properties of the jet are determined from its constituent particles. The typical recombination scheme is to sum all the four-vectors of the constituent to determine the jet four-vector, and this is what is done in the jet algorithms commonly used in ATLAS analyses. The jets used in this analysis are reconstructed with the anti- k_t algorithm [94], which belongs to a wider category of so-called *sequential recombination* algorithms. These sequential recombination algorithms iteratively combine constituents starting from the closest pairs defined by some inter particle distance metric. The anti- k_t algorithm belongs to the family of generalized- k_t algorithms which use similar distance metrics defined as follows [95]:

1. Start with a list of particles in the event, which each have an associated four momentum.
2. For each pair of particles i, j work out the k_t distance

$$d_{ij} = \min(p_{T,i}^{2p}, p_{T,j}^{2p}) \frac{\Delta R_{ij}^2}{R^2}, \quad (4.1)$$

where $\Delta R_{ij} = \sqrt{(y_i - y_j)^2 + (\phi_i - \phi_j)^2}$ is the distance in the $y - \phi$ plane and p is a parameter. Also work out the *beam distance* for each particle i , $d_{iB} = p_{T,i}^{2p}$.

3. Find the minimum of all the d_{ij} and d_{iB} . If the minimum is a d_{ij} , merge particles i and j into a new single particle with four momentum equal to their

sum. If instead a d_{iB} is the minimum, particle i is declared a final jet and removed from the list of particles.

4. Repeat steps 2 and 3 until the list is empty and a set of final jets has been determined.

The behaviour of the generalized- k_t algorithm changes based on the choice of the parameter p . The original k_t algorithm [96, 97] corresponds to $p = 1$, in this case soft emissions with small p_T and therefore small distance will recombine early in the process. The k_t algorithm can be thought of as reversing the perturbative QCD process of successive splittings by starting from the softest emissions and recombining them iteratively with harder and harder constituents. This method has the downside of being sensitive to soft radiation in the event, originating for example due to detector noise or pileup. Alternatively setting $p = -1$ results in the anti- k_t algorithm, this choice favours hard particles and they will be recombined first. The hard jet will grow combining soft particles until it reaches a size of R centred around the original hard emission. The resulting jets will be relatively insensitive to soft radiation, this is the main feature that motivates using the anti- k_t algorithm as the default for ATLAS analyses.

The jets used in this analysis are reconstructed with the anti- k_t algorithm with radius $R = 0.4$ as implemented in the FASTJET [95] software package. The inputs to the jet algorithm from the calorimeter are so-called *topo-clusters*. These are built by clustering calorimeter cells (channels) according to a nearest-neighbour algorithm [98] based on their signal significance relative to the average expected background noise in the cell. Cells with significant signal become seeds and neighbouring cells are added if they satisfy a signal threshold.

The energy of a topo-cluster is defined at the electromagnetic (EM) scale, the energy measured by the calorimeter from EM showers. An event-by-event *origin correction* is applied to the topo-clusters to account for the position of the hard scatter vertex in that particular event. Jets reconstructed from these origin corrected topo-clusters are so-called *EMTopo jets*. EMtopo jets are used exclusively in this analysis, this was the primary jet definition used in ATLAS physics during Run II, the particle flow (*PFlow*) [99] jet definition has since become the default.

EMTopo jets use only the calorimeter clusters as inputs to the jet finding algorithm, where as PFlow uses so called *particle flow objects* (PFOs). PFOs can be either charged (cPFO) or neutral (nPFO), corresponding to the charged and neutral hadrons that make up jets. The tracking information is used to intelligently assign the calorimeter clusters to the appropriate charged and neutral PFOs. The PFlow algorithm achieves better energy resolution for $p_T < 100$ GeV [99].

To account for the response of the calorimeter, EMtopo jets are calibrated by applying energy and pseudorapidity dependent correction factors derived from simulation in the “EM+JES” scheme [100]. Jets are further calibrated using data-driven corrections derived from the p_T balance of jets measured in γ +jet, Z+jet and multijet topologies. All jets are required to have calibrated $p_T > 25$ GeV and rapidity $|y| < 4.4$. To remove most of the jets originating from different pp collisions in the same bunch crossing, the jet vertex tagger (JVT) is used. In particular, the medium JVT criteria is applied to jets with $|\eta| < 2.5$ and $p_T < 120$ GeV [101]. Jets must also fulfil the “loose” jet cleaning criteria [102] that removes jets originating from non-collision background processes. Events containing at least one jet that fails the jet cleaning criteria are vetoed.

Electrons deposit a significant amount of energy in the calorimeters, and will therefore produce one or several topo-clusters. Clusters from an electron will always appear as the dominant part of a jet, so before selecting a jet it is important to ensure that it does not coincide with a good electron. An overlap removal (OR) procedure is applied to ensure that the selected objects are distinct and not double counted, it is applied in two stages. The first stage is applied to jets that pass a minimal pre-selection ($p_T > 25 \text{ GeV}$, $|y| < 4.4$), a jet is removed if it is within a radius $R < 0.2$ of a good electron (one that passes all selection criteria). This stage of the OR is applied to pre-selected jets because the efficiency scale factor of the JVT criteria can only be calculated after OR but before the application of the JVT criteria that occurs in the final jet selection. The second stage of OR is after full object selection. Any lepton that is within a radius $R < 0.4$ of a selected jet is removed, this is essentially an additional isolation requirement on the selected leptons.

4.4 Corrections to the simulation

Corrections are applied to the MC simulation in order to improve the modelling better match the observed conditions in the ATLAS data. When a distribution is filled from simulated events, instead of filling the appropriate bin with +1 for each event as is done with real data, an event specific weight is used instead. To start, the MC event generator provides a generator weight for each event. The generator weight is determined from a variety of factors, such as to combine LO and NLO generation or to account for the efficiency of event filters used in generation. As part of the analysis procedure, additional weights are combined with the generator

weights to correct the simulation to better match the observed data.

The ATLAS data and MC simulations have different distributions of the number of pileup interactions (μ) per event. *Pileup reweighting* applies event-by-event weights that correct for these differences [103].

Other corrections to the simulation are applied via efficiency *scale factors* (SFs), the ratio of some efficiency measured in data to that in simulation. The SF is determined on a per-object level, such as from the muon reconstruction efficiency, and the SFs are combined to determine an overall event SF that is applied to modify the event weight.

The electron identification, reconstruction and isolation efficiencies are measured with J/ψ and $Z \rightarrow ee$ events using a tag-and-probe method [92]. One electron must pass the tightest requirements (the tag) for the event to be tagged, then the second electron (the probe) can be selected using the requirement that the invariant mass of the di-lepton pair matches the J/ψ or Z boson. This is done to avoid a bias on the measured quantities of the second lepton. The probe electron can then be used to measure the efficiencies. This method is used to calculate electron SF dependent on η and/or p_T .

For muons, an analogous tag and probe method as for electrons is used to measure the identification, reconstruction and isolation efficiencies. This makes use of J/ψ and $Z \rightarrow \mu\mu$ events [93] and scale factors are also parametrised in terms of η and/or p_T .

The efficiency of the applied dilepton triggers are also inconsistent between data and simulation, thus a trigger scale factor is required that accounts for this difference. The dielectron triggers applied are symmetric, requiring two electrons

each with a given p_T , as such the electrons are effectively identical in terms of calculating the trigger scale factor. In contrast the dimuon triggers place different p_T requirements on each muon and we cannot know which of the reconstructed muons satisfied each requirement. Thus the muon trigger efficiency accounts for the relative efficiencies of each reconstructed muon to have satisfied each stage of the trigger.

Chapter 5

Analysis

For the first time, thanks to the large dataset collected during Run II of the LHC, it is now feasible to measure the EW- Zjj cross section differentially. The first such measurements are the main result presented in this thesis. First, the event selection selection applied to the full Run II dataset is described; the aim of which is to sufficiently enhance the EW- Zjj signal process. Signal and control regions will also be defined; these regions are used in the statistical method that corrects the modelling of the strong- Zjj component that is described in Section 6.1.

5.1 Strategy

The most common approach when studying Z boson physics at the LHC is to focus on the $Z \rightarrow ee$ and $Z \rightarrow \mu\mu$ final states. As described in Section 2.4, the branching ratio of Z to charged leptons is only 10.1%, the Z decays also to neutrinos ($\mathcal{B} = 20.0\%$) and to all flavours of quarks, except the top quark ($\mathcal{B} = 69.9\%$). The decays to quarks will result in a final state of only hadronic jets, which is very challenging to study due to the enormous background from dijet production (see

Figure 2.8). Neutrinos are not directly measured by the ATLAS detector, instead their presence is inferred in events with missing transverse energy. This adds a layer of difficulty to the neutrino channels. Lastly, the $Z \rightarrow \tau\tau$ channel will produce τ leptons that decay before entering the active area of the detector ($c\tau \simeq 9 \times 10^{-5}$ m) and the τ decay modes contain at least one neutrino and thus suffer from a similar difficulty as in the neutrino channels.

In contrast to the difficulty of the other channels, the $Z \rightarrow ee$ and $Z \rightarrow \mu\mu$ channels produce leptons that are well measured by the ATLAS detector; so these two channels give the best measurement of the Z boson and are the only channels considered in this analysis. Together these two channels are referred to as $Z \rightarrow \ell\ell$. EW- Zjj events are also expected to contain two high p_T jets in the forward regions of the detector coming from the hard scatter quarks, thus $\ell^+\ell^-jj$ is the expected final state.

As previously stated, the expected $\ell\ell jj$ final state is primarily produced by the Drell-Yan process (strong- Zjj). The rarer EW- Zjj process of interest has sufficiently different event kinematics that selection criteria can be devised to greatly enhance the EW component of the selected events. To briefly recap Section 2.5, the EW- Zjj process tends to produce two jets with significant transverse momentum (p_T) that are well separated in rapidity, *i.e.* one in the forward and backward rapidity regions. The Z decay products will be preferentially produced centrally in the gap region between the two high p_T forward jets. This large rapidity gap is also where we expect to see a lack of additional hadronic activity in the EW- Zjj process and the opposite in the strong- Zjj processes. The strong- Zjj process is more likely to have one or more gap jets because the interacting quarks are colour connected. This

results in a difference in final state jet multiplicity that is exploited to enhance the EW component of selected events. The Z boson is expected to be somewhat boosted ($\langle p_T^{\ell\ell} \rangle \approx 150 \text{ GeV}$), and that it is centrally located with respect to the dijet system. Finally, the typical EW- Zjj event is *balanced* in p_T , such that the Z boson candidate balances the dijet system in the transverse plane. The event selection is described in detail in the next section.

The novel differential measurement of the EW- Zjj cross section is an interesting measurement of the SM and it has important applications. It can be used to compare to different theoretical predictions, which can vary significantly, to see what best describes the observed data. The differential measurements are also a powerful tool in the search for new physics, deviations from the expected SM predictions can be studied.

The cross section is measured as a function of four characteristic variables, namely the dijet invariant mass (m_{jj}), the dijet rapidity separation ($|\Delta y_{jj}|$), the p_T of the Z boson ($p_T^{\ell\ell}$) and the signed azimuthal angle between the two jets ($\Delta\phi_{jj}^{\text{signed}}$). The observables m_{jj} and $|\Delta y_{jj}|$ characterize the dijet system and the two observables are strongly correlated. Both are used to define the selection criteria as their dependence is very different between the strong- Zjj and EW- Zjj processes. The $p_T^{\ell\ell}$ observable is of interest because Z production is the focus of the analysis and the Z can be highly boosted with transverse momentum of $\mathcal{O}(1 \text{ TeV})$. High $p_T^{\ell\ell}$ events are also sensitive to higher order operators introduced to the SM, as was studied with the EFT interpretation. Finally, there is the $\Delta\phi_{jj}^{\text{signed}}$ observable, defined as $\Delta\phi_{jj}^{\text{signed}} = \phi_f - \phi_b$ where the two highest p_T jets are ordered such that $y_f > y_b$. The $\Delta\phi_{jj}^{\text{signed}}$ observable is of interest because of its sensitivity to CP-violating effects

in the weak-boson self interaction [104]. This property makes $\Delta\phi_{jj}^{\text{signed}}$ the most sensitive observable to CP-violating beyond the Standard Model (BSM) effects that were tested with the EFT interpretation (Chapter 7).

The differential cross sections are measured separately for inclusive Zjj and for EW- Zjj only. Inclusive Zjj includes both strong- Zjj and EW- Zjj as well as a relatively small contribution from diboson Zjj . The EW- Zjj only process is measured in a fiducial signal region (SR) with an enhanced component of the EW signal. The inclusive process is measured in the EW signal region as well as in the control regions used in the EW- Zjj signal extraction procedure.

The differential cross section for EW- Zjj production as a function of a variable x is defined in each bin i of the distribution by

$$\frac{d\hat{\sigma}_i^{\text{EW}}}{dx} = \frac{\mathcal{U}\hat{N}_{\text{SR},i}^{\text{EW}}}{\Delta x_i \mathcal{L}} = \frac{\mathcal{U} \left(N_{\text{SR},i}^{\text{data}} - \hat{N}_{\text{SR},i}^{\text{strong}} - \hat{N}_{\text{SR},i}^{\text{non-Z}} \right)}{\Delta x_i \mathcal{L}}, \quad (5.1)$$

where $\mathcal{L}$ is the integrated luminosity of the dataset, Δx_i is the width of the bin and SR represents the signal region (detailed in Section 5.2). The $\hat{N}_{\text{SR},i}^{\text{EW}}$ term is the estimated number of EW Zjj events, defined as the data minus the non-EW background. $\mathcal{U}$ represents a set of corrections for the detector effects (Section 6.7), the procedure of evaluating and applying these corrections is typically referred to as *unfolding*. The small non- Zjj contribution ($\hat{N}_{\text{SR},i}^{\text{non-Z}}$) comes primarily from diboson and $t\bar{t}$ processes and is taken from MC predictions normalized to the best available cross sections. The largest background contribution comes from strong Zjj production ($\hat{N}_{\text{SR},i}^{\text{strong}}$), which is estimated using data-driven approach described in Section 6.1.

5.2 Event selection and fiducial regions

The event selection is designed to enhance the EW signal while rejecting various backgrounds, with the most significant being strong- Zjj . The selection is similar to the selection used in previous EW- Zjj analyses [5][7].

The data are collected using a dilepton trigger (Table 4.3) and are required to fulfil the data quality criteria listed in Section 4.2. The first stage of the event selection requires exactly two reconstructed same flavour opposite sign leptons with $m_{\ell\ell} \in (81 \text{ GeV}, 111 \text{ GeV})$. The selected leptons must satisfy the electron or muon object selection described in Sections 4.3.2 and 4.3.3, respectively. This *dilepton* selection is designed to select events containing exactly one Z boson candidate. The selected leptons are labelled ℓ_1 and ℓ_2 , ordered such that $p_T^{\ell_1} > p_T^{\ell_2}$. These will either be muons or electrons, separating the events into 2 channels that are compared against one another at this level of selection to ensure consistency. The channels are then combined and further selection criteria are applied identically to events from either channel.

The data to MC agreement after dilepton selection is shown in Figures 5.1–5.4. These comparisons serve as the first checks that the normalization and corrections to the simulation have been applied correctly. Good agreement is observed across the full range of p_T and η of the leptons. The NLO SHERPA prediction describes the data somewhat better than the LO MG5+PY8 prediction overall where some discrepancy is observed at high lepton p_T . The agreement in terms of the average number of interactions per bunch crossing, μ , and the number of reconstructed primary vertices, N_{PV} , in Figure 5.4 indicates that the pileup reweighting (Section 4.4) has been correctly applied and that the three MC campaigns (see Section 4.1.2) have

been combined correctly.

Next, the *dijet* selection requires two or more jets where the two highest p_T jets fulfil $p_T^{j1} > 85 \text{ GeV}$ and $p_T^{j2} > 80 \text{ GeV}$. This selection criteria is designed to select events that will fit the VBF Z topology where we expect two high p_T jets in the forward and backward rapidity region, respectively. As is shown in Figure 5.5, EW- Zjj has a significantly harder jet p_T spectrum compared to strong- Zjj and the other backgrounds. The leading two jets are denoted $j1$ and $j2$ and their four momentum sum define the dijet system, denoted jj . Any jet that has rapidity in the rapidity interval of the two leading jets (*i.e.* between y_{j1} and y_{j2}) is defined as a gap jet. The multiplicity of such jets, $N_{\text{jets}}^{\text{gap}}$, is a key observable of the analysis. Note that gap jets need to fulfil the jet object definition outlined in Section 4.3.4 ($p_T > 25 \text{ GeV}$, $|y| < 4.4$).

The data to MC agreement after the dijet selection is applied is quite poor, the requirement of two high p_T jets with a large rapidity separation defines a phase space where the MC modelling is challenging. This modelling difficulty has been observed previously [5, 105], particularly the SHERPA event generator tends to over-predict the observed data for jet observables in the forward region. The data to MC agreement after dilepton selection is shown in Figures 5.6–5.8. At high values of m_{jj} ($> 1.5 \text{ TeV}$) and $|\Delta y_{jj}|$ (> 5) the SHERPA strong- Zjj prediction is two - three times the observed data. There is also shows poor agreement in both the jet and gap jet multiplicities. The SHERPA and MG5+PY8 predictions have very different shapes with respect to one another and to the data.

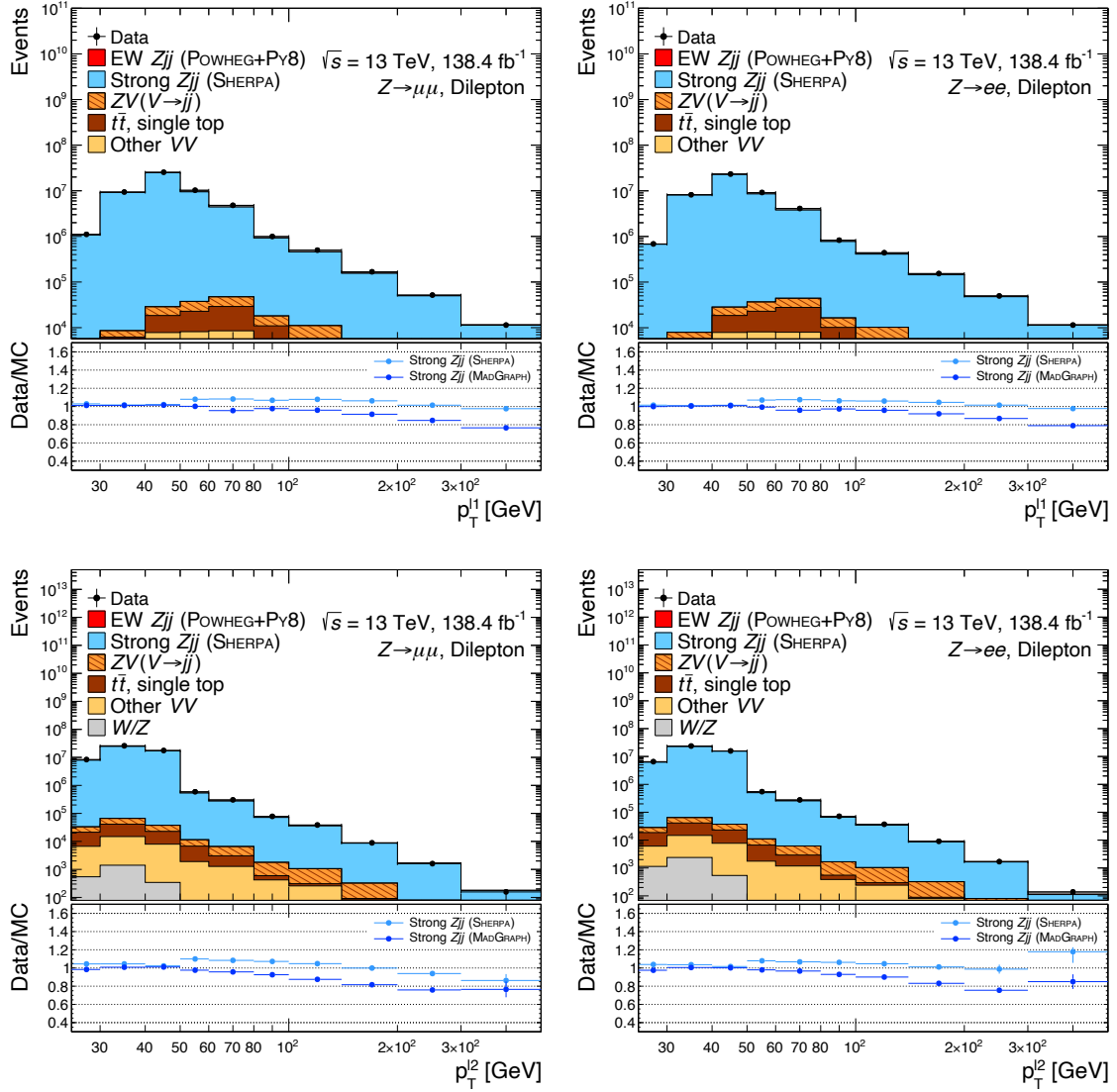


FIGURE 5.1: Observed and predicted distributions of p_T^{l1} (top) and p_T^{l2} (bottom) for the $Z \rightarrow \mu\mu$ (left) and $Z \rightarrow ee$ (right) channels after the dilepton selection.

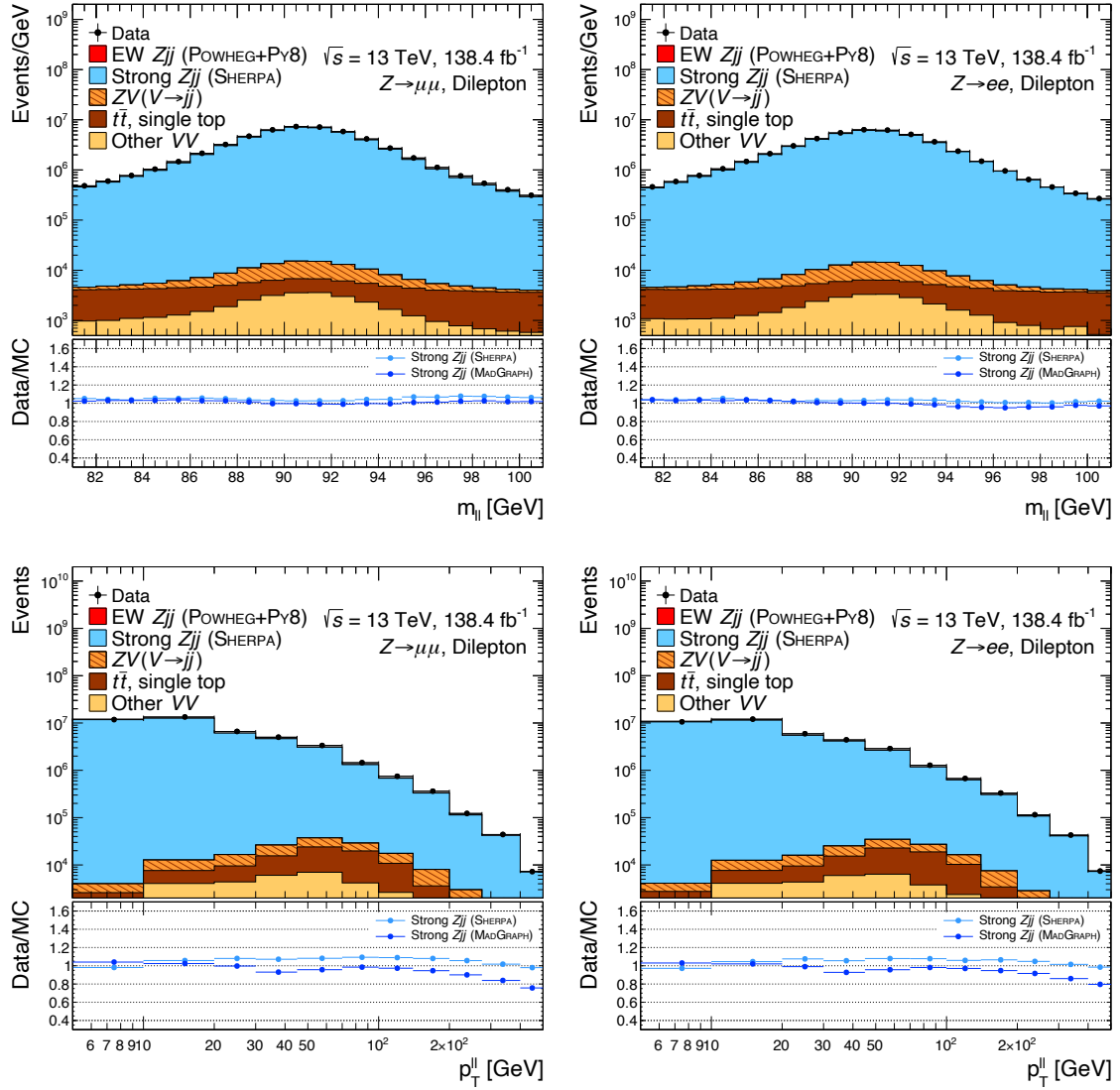


FIGURE 5.2: Observed and predicted distributions of $m_{\ell\ell}$ (top) and $p_T^{\ell\ell}$ (bottom) for the $Z \rightarrow \mu\mu$ (left) and $Z \rightarrow ee$ (right) channels after the dilepton selection.

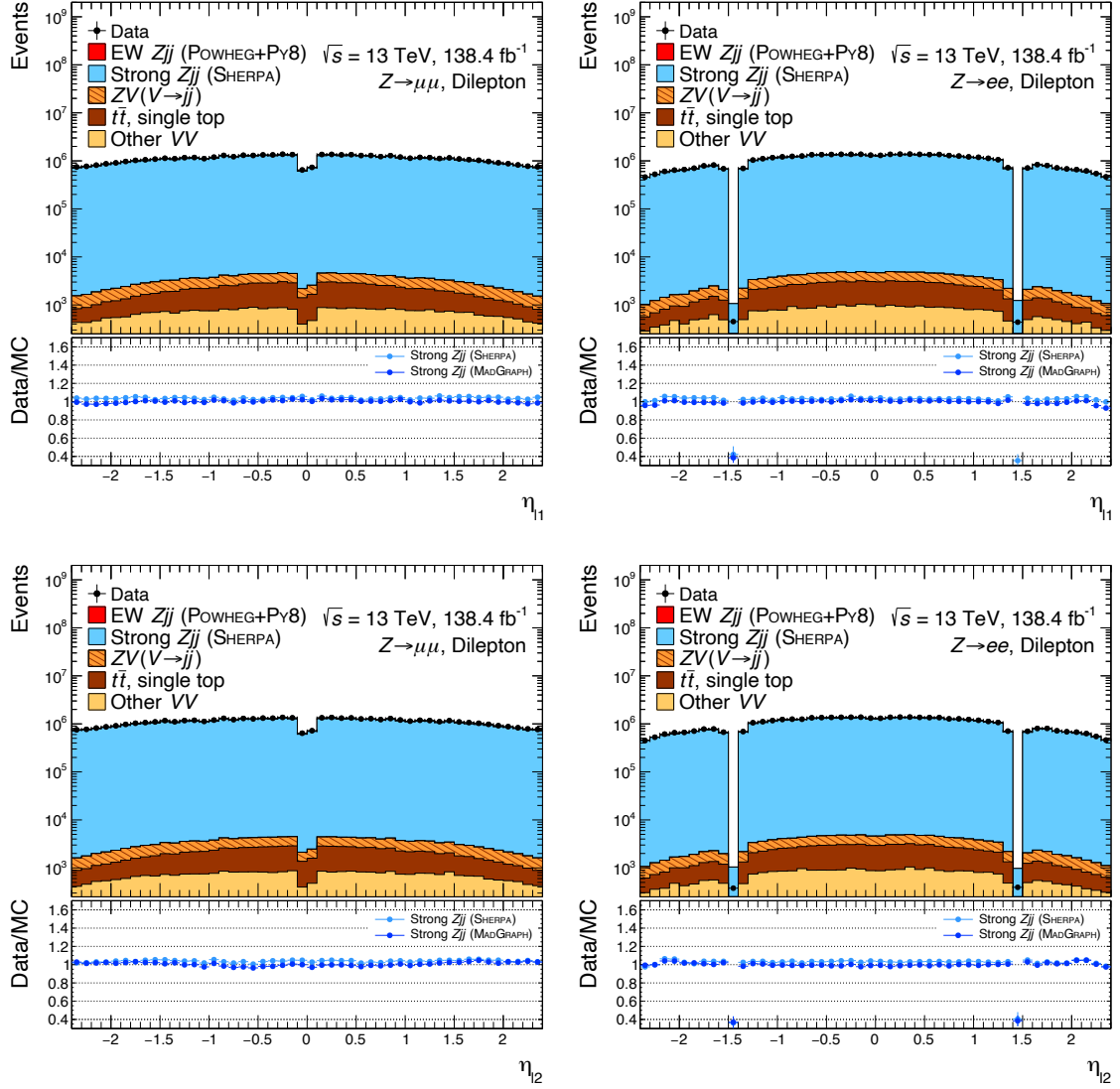


FIGURE 5.3: Observed and predicted distributions of $\eta^{\ell 1}$ (top) and $\eta^{\ell 2}$ (bottom) for the $Z \rightarrow \mu\mu$ (left) and $Z \rightarrow ee$ (right) channels after the dilepton selection.

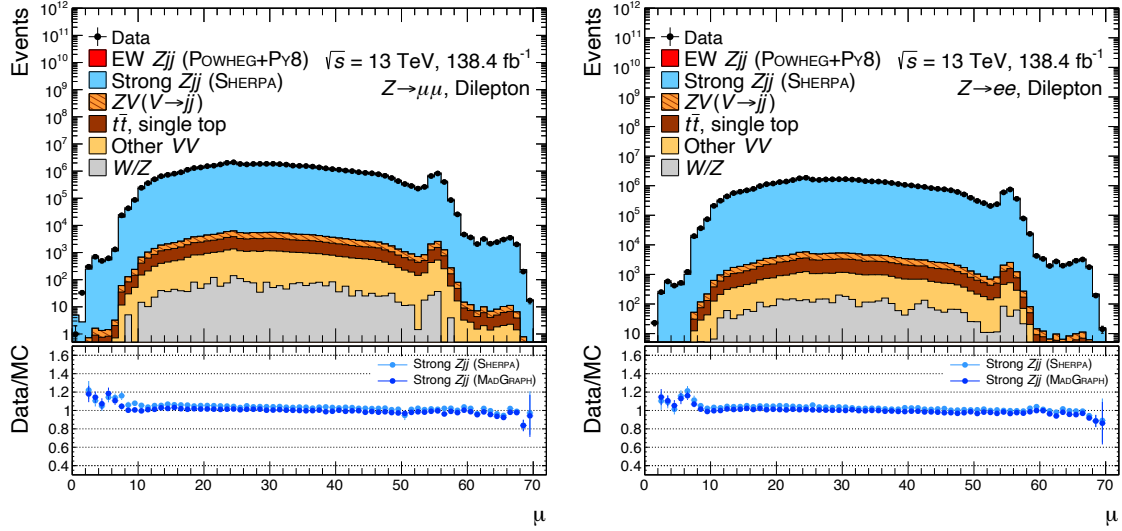


FIGURE 5.4: Distribution of the average interactions per bunch crossing (μ) for data and simulation in the $Z \rightarrow \mu\mu$ channel (left) and the $Z \rightarrow ee$ channel (right). The discrepancies observed in the ratio at the edge values of the distribution are concentrated at values of μ where a very small fraction of the total events were observed. These small discrepancies will not impact the measurements of EW- Zjj since the process is relatively insensitive to pileup.

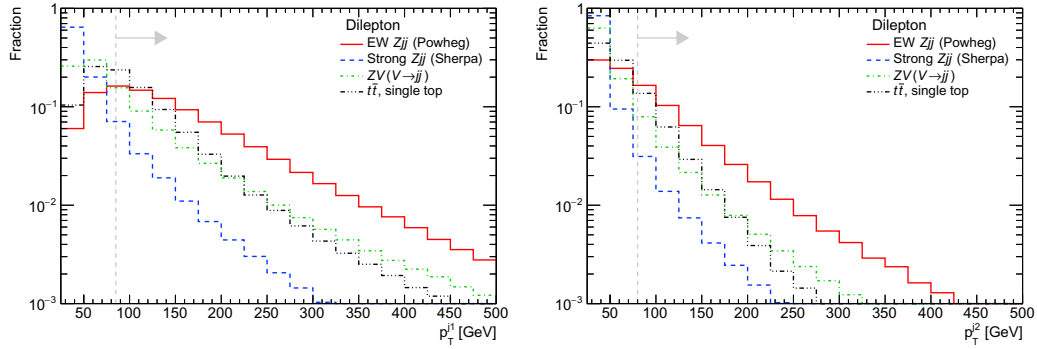


FIGURE 5.5: Shape comparison of jet p_T from EW- Zjj , strong- Zjj , Dilepton Zjj and top events. The jet observables p_T^{j1} (left) and p_T^{j2} (right) are shown after dilepton selection with the dijet selection criteria value depicted in gray for each observable, the arrows indicate the accepted region.

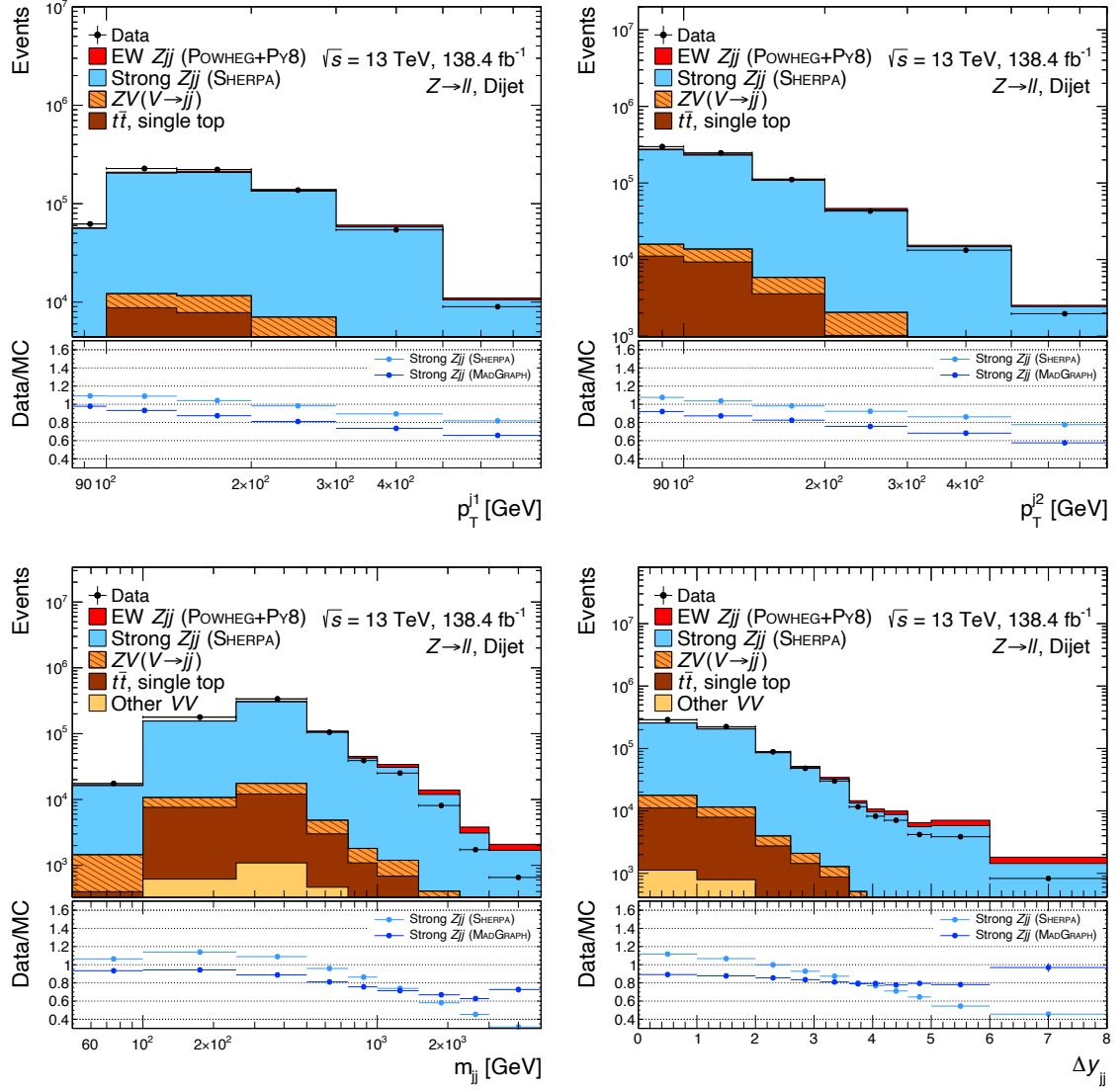


FIGURE 5.6: Observed and predicted distributions of p_T^{j1} , p_T^{j2} (top), m_{jj} and $|\Delta y_{jj}|$ (bottom) after the dijet selection.

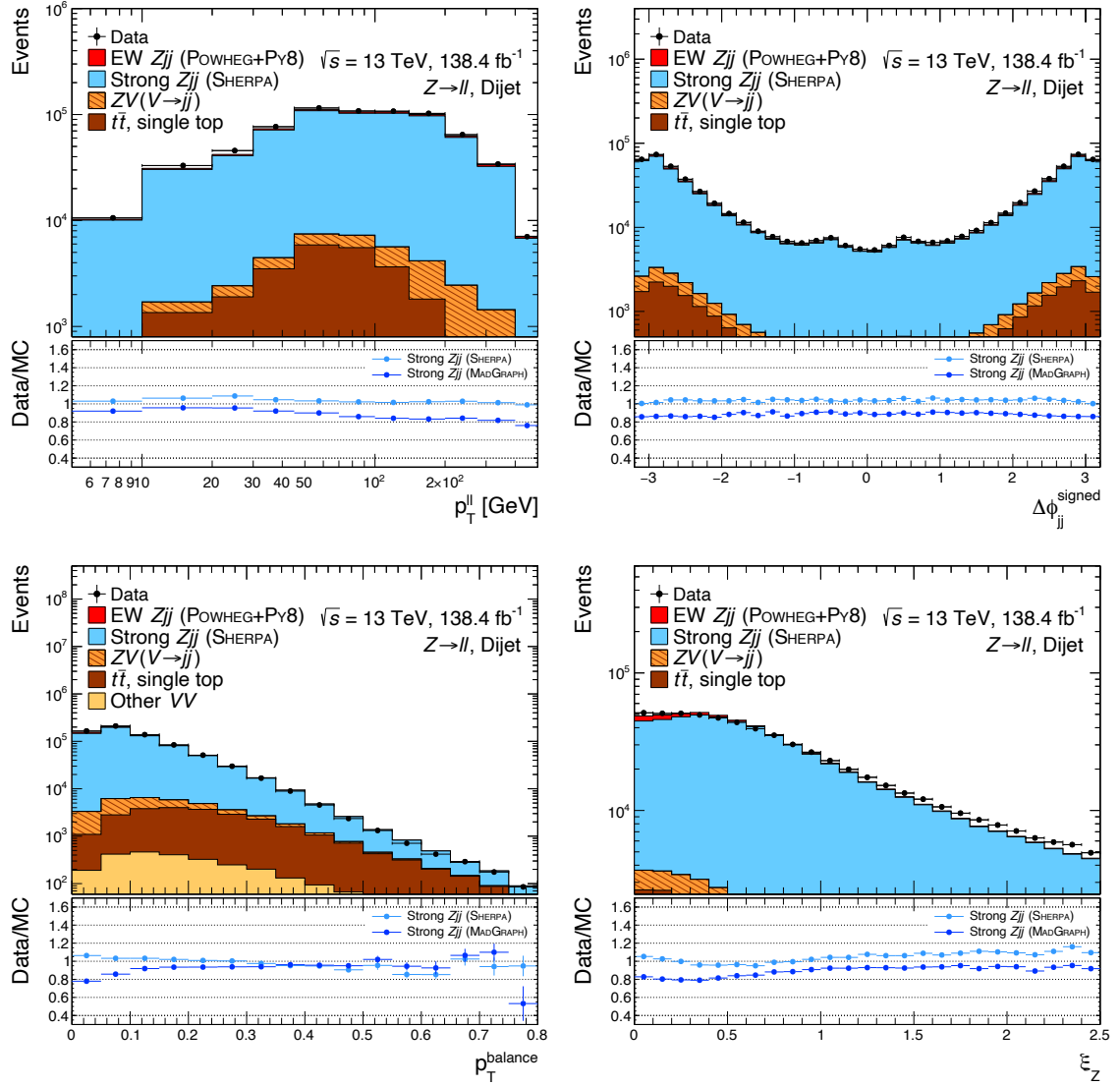


FIGURE 5.7: Observed and predicted distributions of $p_T^{\ell\ell}$, $\Delta\phi_{jj}^{\text{signed}}$ (top), p_T^{bal} and ζ_Z (bottom) after the dijet selection.

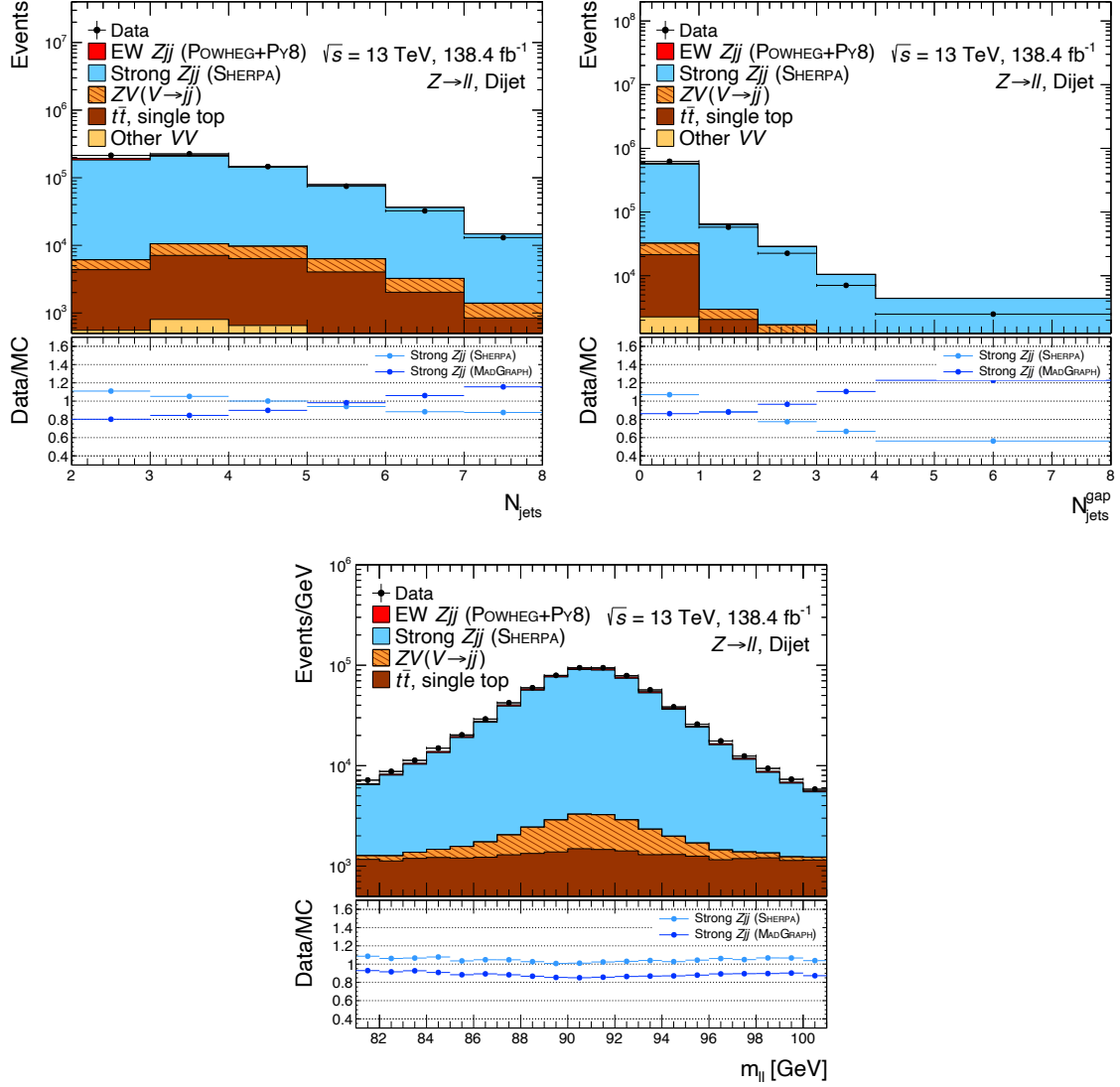


FIGURE 5.8: Observed and predicted distributions of N_{jets} , $N_{\text{jets}}^{\text{gap}}$ (top) and $m_{\ell\ell}$ (bottom) after the dijet selection.

Following dijet selection, the next stage of event selection requires two specialized observables. First, the p_T balance observable is defined as

$$p_T^{\text{bal}} \equiv |\sum \vec{p}_T| / \sum p_T, \quad (5.2)$$

where the sum runs over the two leptons and the leading two jets, and the leading gap jet, if such a jet exists. Next, the Z centrality observable ξ_Z is defined as

$$\xi_Z = \frac{|y_{\ell\ell} - (y_{j1} + y_{j2})/2|}{|\Delta y_{jj}|} \simeq \frac{y_{\ell\ell} - y_{jj}}{|\Delta y_{jj}|}, \quad (5.3)$$

i.e. the difference in rapidity of the dilepton and dijet systems over the rapidity separation of the two leading jets.

The remaining selection criteria are applied to events passing dijet selection to establish the *VBF topology* selection. The first two criteria are $m_{jj} > 1 \text{ TeV}$ and $|\Delta y_{jj}| > 2.0$. Selecting events with a dijet system with high invariant mass and large rapidity separation serves to enhance the EW- Zjj signal component. Figure 5.9 shows how EW- Zjj tends to occur at higher m_{jj} and $|\Delta y_{jj}|$ compared to the other processes. These two variables, and particularly m_{jj} , are the most powerful observables for distinguishing between strong and EW Zjj .

The $|\Delta y_{jj}| > 2.0$ requirement is also designed to minimize the quantum mechanical interference between strong- Zjj and EW- Zjj . The analysis strategy involves subtracting the strong Zjj component from the observed data to extract the EW signal. This subtraction is performed assuming there is no significant interference between the strong- Zjj and EW- Zjj processes. This interference was studied with the MADGRAPH event generator by applying the event selection at particle-level

to two samples, one producing EW- Zjj events and the other containing the interference contribution. As shown in Figure 5.10, the interference has the largest impact for $|\Delta y_{jj}| < 1.0$. At $|\Delta y_{jj}| = 1$ the fractional interference is unity. However, when $|\Delta y_{jj}| > 2.0$ the interference contribution is always $< 10\%$ and continues to decrease with increasing $|\Delta y_{jj}|$. The fractional interference is smaller for the other measurement variables (m_{jj} , $p_T^{\ell\ell}$ and $\Delta\phi_{jj}^{\text{signed}}$) for which the impact ranges from 0–7% of the EW- Zjj contribution. The impact of the interference on the final measurement is considered as a systematic uncertainty (see Section 6.4.2).

Next is the requirement that the Z boson is somewhat boosted, $p_T^{\ell\ell} > 20 \text{ GeV}$. The main reason for this requirement is to suppress events where a Z boson is produced with no hard scatter jets and containing two high p_T forward pileup jets. These forwards jets are common from pileup dijet events and cannot be suppressed by the JVT requirement as they occur outside the tracker. Requiring $p_T^{\ell\ell} > 20 \text{ GeV}$ increases the likelihood that the observed jets are produced from the hard scatter. The Z p_T spectrum is also slightly harder in EW- Zjj compared to strong- Zjj .

The event is required to be balanced in p_T , this requirement of $p_T^{\text{bal}} < 0.15$ will particularly suppress $t\bar{t}$ and single top events as the neutrinos produced in the decay of top quarks will tend to produce a dijet and dilepton system that are imbalanced relative to those of the EW- Zjj signal. This selection will also suppress WW events since leptonic W boson decay will also produce a neutrino. Low Z centrality ($\xi_Z < 1.0$) is required because for EW- Zjj the $y_{\ell\ell}$ and y_{jj} quantities will be similar and thus Z centrality will be small from Equation (5.3). The shape of each selection variable for EW- Zjj and the three most important background processes (strong- Zjj , diboson Zjj and top) are shown in Figure 5.9. The selection

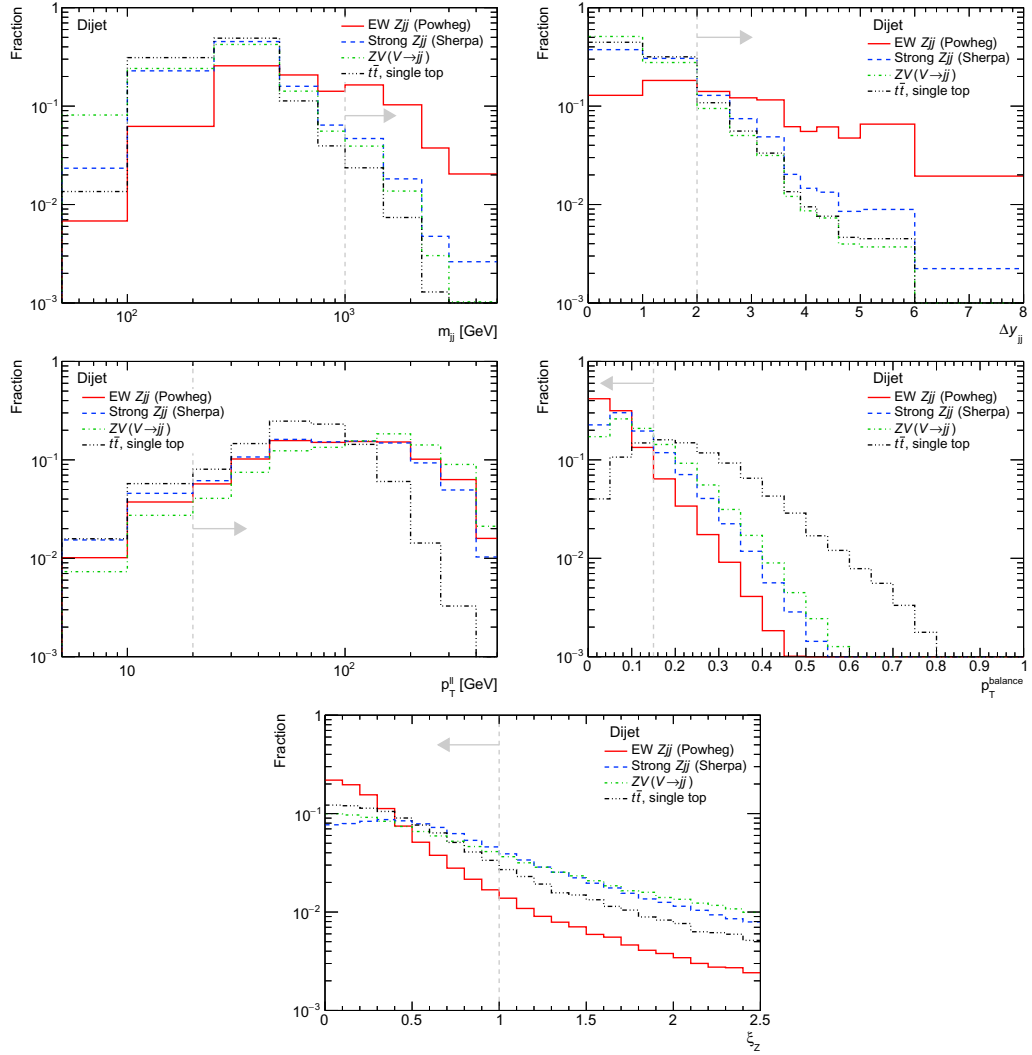


FIGURE 5.9: Shape comparison of EW- Zjj , strong- Zjj , Diboson Zjj and top events in key selection variables after dijet selection. The selection criteria on these variables are applied after dijet selection to define the VBF topology region. The value applied to each variable is shown in gray, the arrows indicate the accepted region (see Table 5.1).

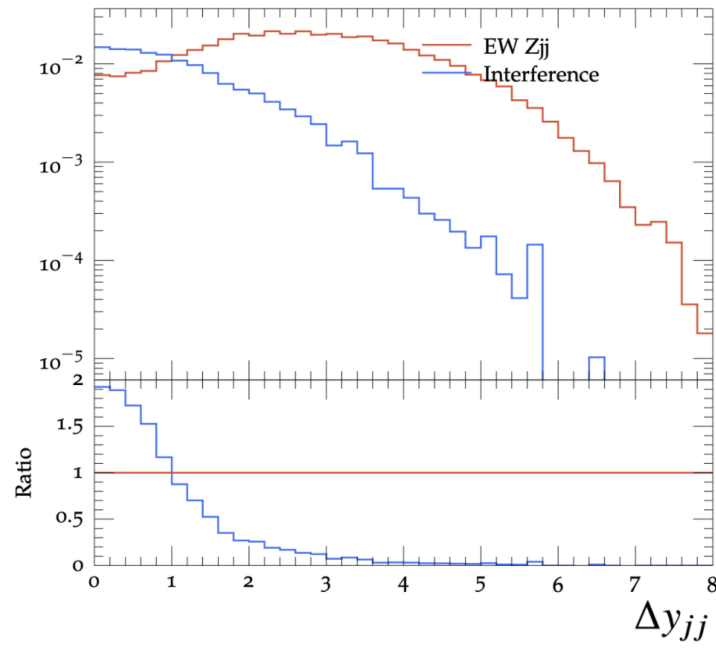


FIGURE 5.10: Differential distribution of EW- Zjj and the interference between strong- Zjj and EW- Zjj as a function $|\Delta y_{jj}|$ after applying the dijet event selection.

is summarized in Table 5.1.

Event selection	Description
Trig+DQ	Dilepton trigger, lumiblock flagged good by GRL, data quality flags
Dilepton	Two opposite sign leptons e^+e^- or $\mu^+\mu^-$ with $m_{\ell\ell} \in (81 \text{ GeV}, 111 \text{ GeV})$
Dijet	At least two jets where the leading two have $p_T^{j1} > 85 \text{ GeV}$ and $p_T^{j2} > 80 \text{ GeV}$
VBF topology	$m_{jj} > 1 \text{ TeV}$, $ \Delta y_{jj} > 2.0$, $p_T^{\ell\ell} > 20 \text{ GeV}$, $p_T^{\text{bal}} < 0.15$, $\xi_Z < 1.0$

TABLE 5.1: Overview of the event selection. The criteria are applied cumulatively.

The impact of the event selection on the data and MC sample event selection efficiency is shown in Tables 5.2 and 5.3 for the $Z \rightarrow ee$ and $Z \rightarrow \mu\mu$ channels, respectively. The MC simulation has been normalized to match the luminosity of the Run-II dataset. The offset in the event total between data and MC in the first two columns is due to the presence of QCD background ($pp \rightarrow X$) in the data. These QCD events are very efficiently rejected by the selection, particularly the lepton quality and isolation requirements. The m_{jj} requirement in these tables is $m_{jj} > 250 \text{ GeV}$, significantly lower than the final value of $m_{jj} > 1 \text{ TeV}$ that partially defines the VBF topology selection. This lower m_{jj} requirement greatly enhances the strong- Zjj component of the selection, it is used for the measurement of inclusive Zjj as a function of m_{jj} . Since the $m_{jj} > 250 \text{ GeV}$ is applied after the $|\Delta y_{jj}| > 2.0$ requirement and since these are strongly correlated the m_{jj} selection criteria has little effect in this table. The final requirement of $m_{jj} > 1 \text{ TeV}$ does make a significant impact on the event composition of the VBF topology region.

Table 5.4 shows the event yield of each sample in both lepton channels after the full VBF topology selection with $m_{jj} > 1 \text{ TeV}$, comparing the final column of Tables 5.2 and 5.3 shows the significant impact of raising the m_{jj} criteria on the

event yield. After VBF topology selection the composition of the selected MC events is $\sim 15\%$ EW- Zjj . Comparing this to dijet selection where EW- Zjj account for only $\sim 0.05\%$ of event illustrates how the applied selection requirements succeed in enhancing the EW- Zjj signal component.

$Z \rightarrow ee$ channel		Selection							
Sample	None	Trig+DQ	Dilepton	Dijet	$ \Delta y_{jj} $	m_{jj}	$p_T^{\ell\ell}$	p_T^{balance}	ξ_Z
Data 2015–2018	925.3M	96.1M	46.8M	342k	95.9k	95.6k	88.9k	67.0k	63.2k
EW- Zjj (POWHEG+PY8)	106k	48.8k	26.7k	9.9k	6.8k	6.8k	6.5k	5.8k	5.8k
EW- Zjj (SHERPA)	87.4k	35.3k	18.7k	6.2k	3.7k	3.7k	3.5k	3.2k	3.1k
Strong- Zjj (SHERPA)	289.7M	88.6M	47.9M	325k	103k	103k	95.6k	71.3k	66.9k
Strong- Zjj (MG5+PY8)	292.9M	91.6M	49.4M	389k	111k	111k	104k	83.3k	79.2k
$ZV(V \rightarrow jj)$	340k	108k	60.8k	6.8k	1.4k	1.4k	1.4k	906	871
Other VV	14.1M	994k	65.1k	12.1k	2.9k	2.9k	2.7k	785	772
$t\bar{t}$, single top	1.0M	150k	29.8k	1.4k	407	406	390	166	158
$W(\rightarrow \ell\nu)+\text{jets}, Z(\rightarrow \tau\tau)+\text{jets}$	1.9B	611k	4.4k	15	2	2	2	2	2

TABLE 5.2: $Z \rightarrow ee$ channel data event yields and predicted yields from MC at various stages of the event selection. The values of the selection criteria applied in the last five columns are those listed in Table 5.1 except for the m_{jj} requirement, which is $m_{jj} > 250$ GeV.

$Z \rightarrow \mu\mu$ channel		Selection							
Sample	None	Trig+DQ	Dilepton	Dijet	$ \Delta y_{jj} $	m_{jj}	$p_T^{\ell\ell}$	p_T^{balance}	ξ_Z
Data 2015–2018	925.3M	349.7M	52.6M	372k	106k	106k	98.5k	74.5k	69.5k
EW- Zjj (POWHEG+PY8)	106k	68.5k	29.8k	10.6k	7.3k	7.3k	6.9k	6.2k	6.2k
EW- Zjj (SHERPA)	87.9k	52.9k	21.2k	6.7k	4.1k	4.1k	3.8k	3.4k	3.4k
Strong- Zjj (SHERPA)	289.4M	121.0M	55.0M	352k	114k	114k	106k	78.4k	73.2k
Strong- Zjj (MG5+PY8)	291.4M	125.5M	56.8M	425k	124k	124k	116k	92.5k	87.5k
$ZV(V \rightarrow jj)$	340k	154k	68.0k	7.1k	1.5k	1.5k	1.4k	937	898
Other VV	14.1M	4.8M	68.9k	12.2k	2.9k	2.9k	2.7k	825	817
$t\bar{t}$, single top	1.0M	285k	33.8k	1.5k	435	434	418	176	166
$W(\rightarrow \ell\nu)+\text{jets}, Z(\rightarrow \tau\tau)+\text{jets}$	1.9B	7.9M	2.5k	20	0	0	0	0	0

TABLE 5.3: $Z \rightarrow \mu\mu$ channel data event yields and predicted yields from MC at various stages of the event selection. The values of the selection criteria applied in the last five columns are those listed in Table 5.1 except for the m_{jj} requirement, which is $m_{jj} > 250$ GeV.

Sample	$Z \rightarrow ee$				$Z \rightarrow \mu\mu$			
Data	10870				12125			
EW- Zjj (POWHEG+PY8)	2670	$\pm$ 120	$\pm$ 280		2740	$\pm$ 120	$\pm$ 290	
EW- Zjj (SHERPA)	1280	$\pm$ 60	$\pm$ 140		1350	$\pm$ 60	$\pm$ 150	
EW- Zjj (HERWIG+VBFNLO')	2290	$\pm$ 100	$\pm$ 210		2350	$\pm$ 100	$\pm$ 220	
Strong- Zjj (SHERPA)	13500	$\pm$ 600	$\pm$ 4500		15100	$\pm$ 600	$\pm$ 5000	
Strong- Zjj (MG5+PY8)	13140	$\pm$ 480	$\pm$ N/A		14840	$\pm$ 540	$\pm$ N/A	
Strong- Zjj (MG5_NLO+PY8')	8800	$\pm$ 300	$\pm$ 1000		10000	$\pm$ 400	$\pm$ 1200	
$ZV(V \rightarrow jj)$	179	$\pm$ 8	$\pm$ 6		178	$\pm$ 8	$\pm$ 6	
Other VV	45	$\pm$ 2	$\pm$ 2		45	$\pm$ 2	$\pm$ 2	
$t\bar{t}$, single top	92	$\pm$ 8	$\pm$ 6		98	$\pm$ 8	$\pm$ 6	
$W(\rightarrow \ell\nu)$ +jets, $Z(\rightarrow \tau\tau)$ +jets	negligible				negligible			

TABLE 5.4: Observed and expected event yields in the $Z \rightarrow ee$ and $Z \rightarrow \mu\mu$ channels following VBF topology selection. The first (second) uncertainty quoted for each generator is the experimental (theoretical) systematic uncertainty. Theoretical uncertainties are calculated for all predictions except for MG5+PY8 strong- Zjj , which is denoted 'N/A' in the table. $W(\rightarrow \ell\nu)$ +jets, $Z(\rightarrow \tau\tau)$ +jets are considered 'negligible' since a total of two events were selected across both processes and channels.

Figures 5.11 and 5.12 compare the four measurement variables, as well as the two variables that define the signal and control regions (ξ_Z and $N_{\text{jets}}^{\text{gap}}$) after VBF topology selection. Also shown are the ratio of MC to data for each MC sample, this highlights the differences in modelling of the three strong- Zjj samples, particularly in m_{jj} . The SHERPA and MG5+PY8 predictions vastly over and under predict the high m_{jj} ($m_{jj} > 2 \text{ TeV}$) behaviour observed in data, respectively.

5.3 Signal and control regions

After applying the full VBF topology selection the strong- Zjj taken directly from the MC generators significantly over-predicts the data in most bins of the relevant distributions. Thus if Equation (5.1) was applied using the default MC predictions

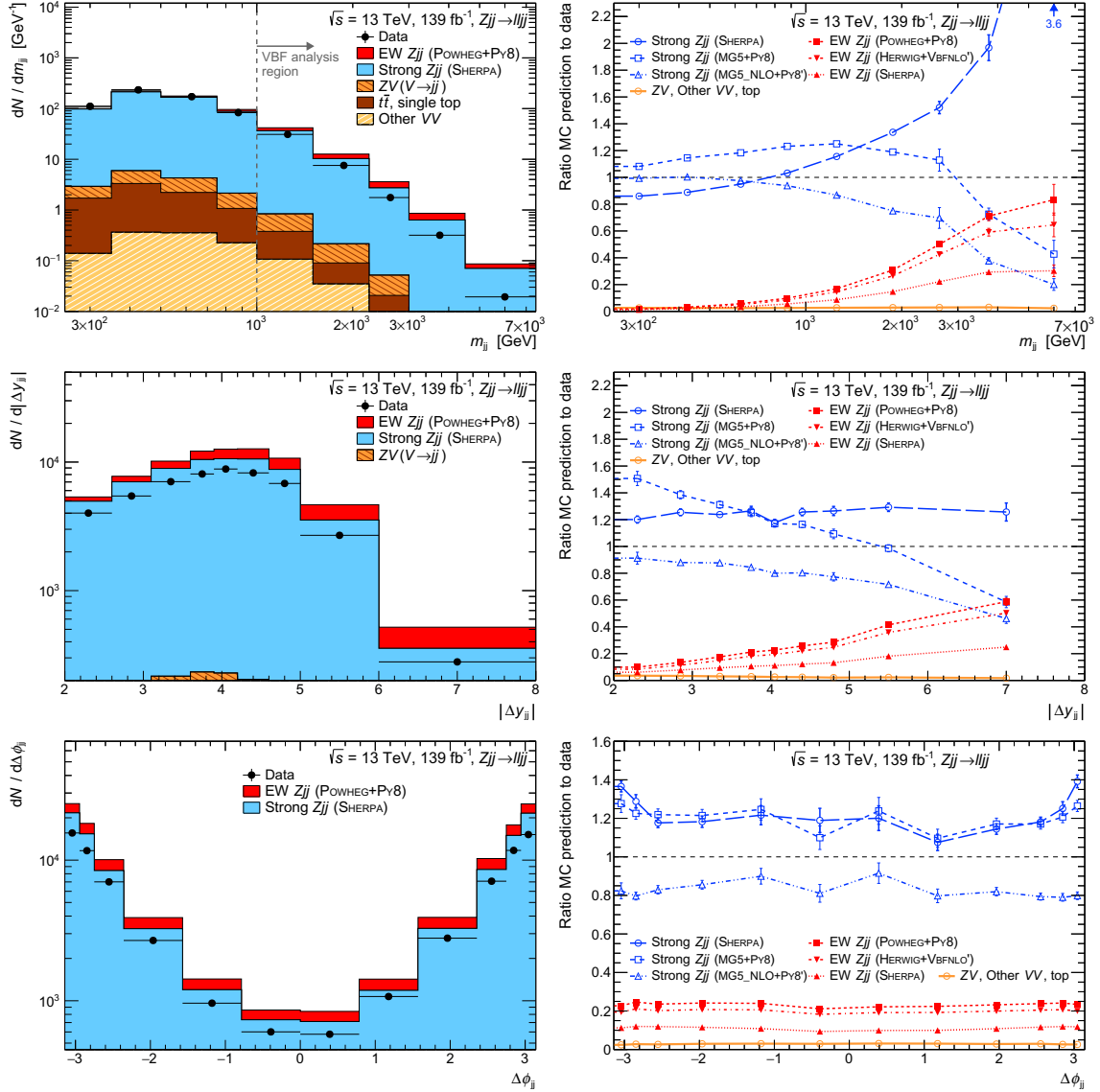


FIGURE 5.11: Observed and predicted distributions of m_{jj} (top), $|\Delta y_{jj}|$ (middle) and $\Delta\phi_{jj}^{\text{signed}}$ (bottom) after VBF topology selection. The data and MC spectra are compared on the left. The ratio of the MC simulation to data for each strong- Zjj and EW- Zjj generator sample is shown on the right, as well as the ratio of the sum of the minor backgrounds. The MG5_NLO+PY8 strong prediction and the POWHEG+PY8 and HERWIG+VBFNLO EW predictions best describe the observed data.

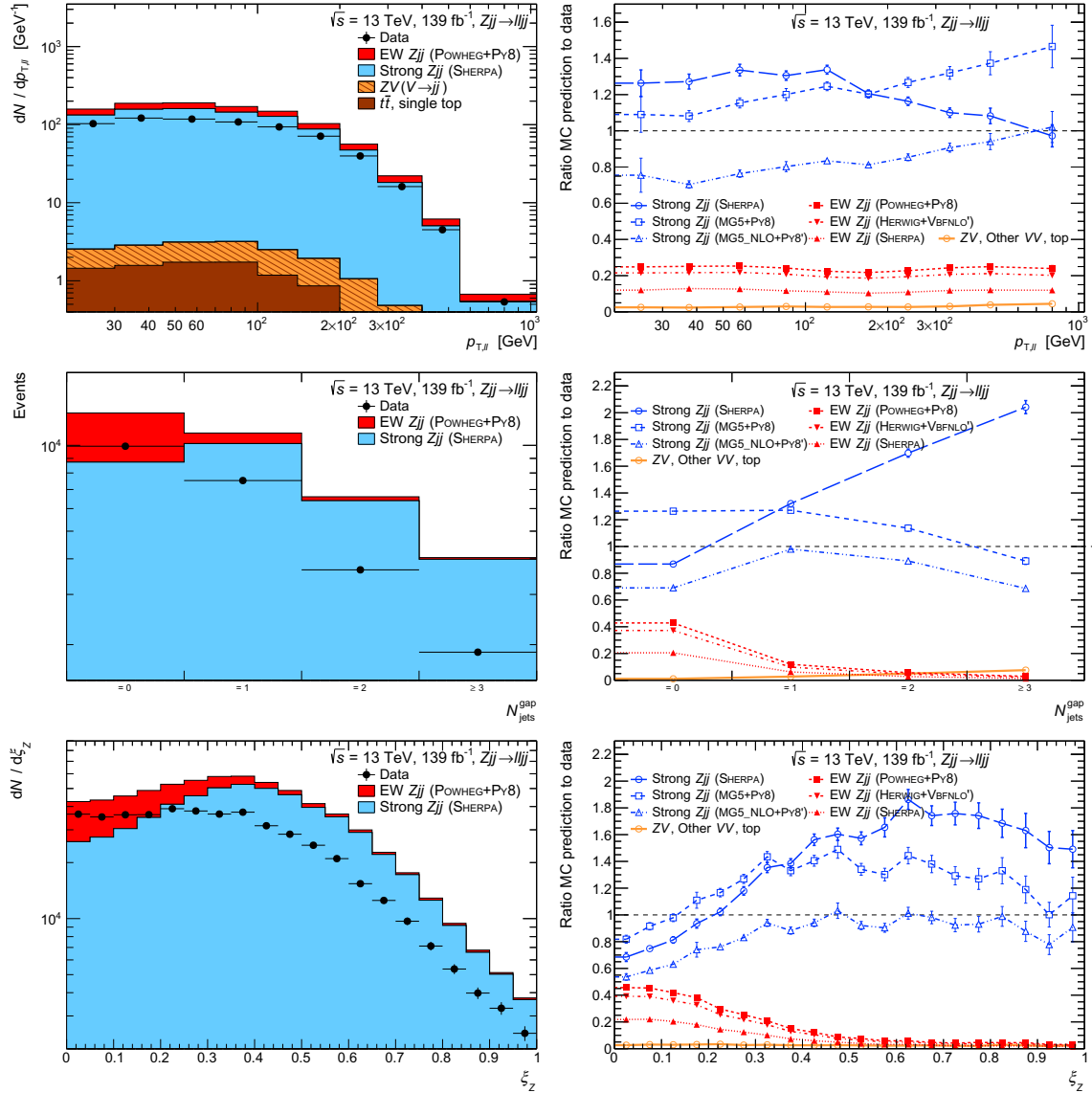


FIGURE 5.12: Observed and predicted distributions of $p_T^{\ell\ell}$ (top), $N_{\text{jets}}^{\text{gap}}$ (middle) and ξ_Z (bottom) after VBF topology selection. The data and MC spectra are compared on the left. The ratio of the MC simulation to data for each strong- Zjj and EW- Zjj generator sample is shown on the right, as well as the ratio of the sum of the minor backgrounds. The combination of the MG5_NLO+PY8 strong prediction and the POWHEG+PY8 or HERWIG+VBFNLO EW predictions best describe the observed data.

the extracted EW- Zjj cross section would be negative. To solve this problem a data driven correction was implemented to correct the strong- Zjj prediction in the EW- Zjj enhanced signal region.

To accomplish this the $N_{\text{jets}}^{\text{gap}}$ and ξ_Z observables are used to subdivide the VBF topology phase space into signal and control regions. These two observables are chosen as they are almost entirely uncorrelated, in data the correlation coefficient ρ was found to be 0.076, defined from the standard deviations and covariance of the two variables as $\rho(X, Y) = \text{cov}(X, Y) / \sigma_X \sigma_Y$. The two-dimensional distribution of $N_{\text{jets}}^{\text{gap}}$ and ξ_Z in the data after VBF topology selection is shown in Figure 5.13. $N_{\text{jets}}^{\text{gap}} = 0$ and $\xi_Z < 0.5$ define the signal region (SR) where EW- Zjj is maximally enhanced. Inverting one or both of these cuts defines 3 control regions (CRa, CRb and CRc) where EW- Zjj is suppressed so these later three control regions contain almost entirely strong- Zjj events. The control regions can thus be used to constrain the strong- Zjj prediction and extrapolate that prediction to determine the strong- Zjj component in the signal region, after which the EW signal can be extracted. This EW signal extraction process is detailed in Section 6.1.

The definition of the signal and control regions is illustrated in Figure 5.14 and the region compositions in terms of EW- Zjj , strong- Zjj and other background are shown in Figure 5.15.

The fiducial definition is very important as it defines the measurement. It is defined at stable particle-level, and the selection criteria closely replicate the ones applied at analysis level. The fiducial selection is applied to simulated events at the particle-level. Leptons are defined at the *dressed* level, as the four-momentum combination of a prompt electron or muon (that do not originate from the decay of

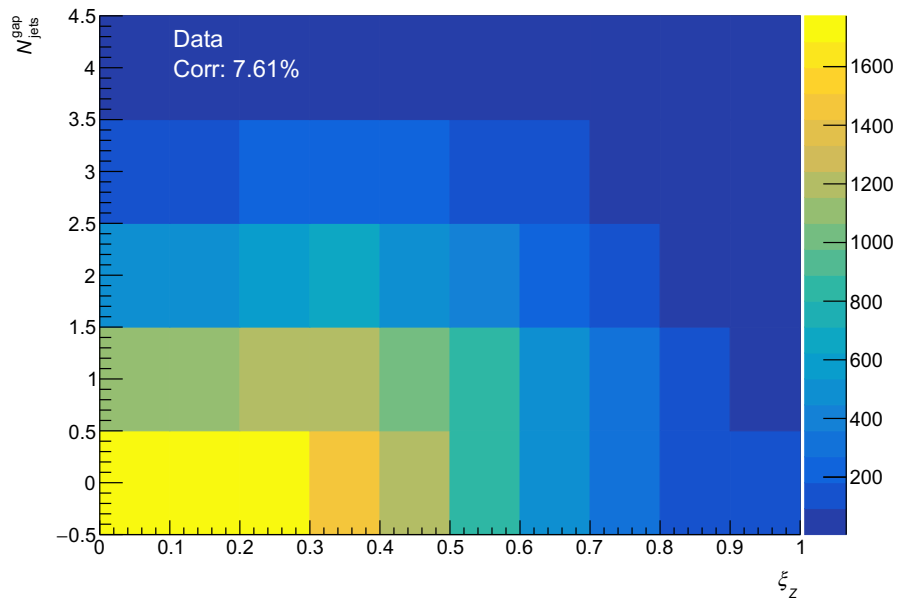


FIGURE 5.13: Two-dimensional distribution of $N_{\text{jets}}^{\text{gap}}$ and ξ_Z after VBF topology selection. The colour indicates the number of data events in each bin. This demonstrates that these two variables are not correlated and therefore are useful for defining the signal and control regions.

a hadron) and all nearby prompt photons within $\Delta R < 0.1$. Leptons are required to have $p_T > 25 \text{ GeV}$ and have the same acceptance requirement as used at the analysis level, *i.e.* muons satisfy $|\eta| < 2.4$ and electrons satisfy $|\eta| < 2.37$ (but exclude the region $1.37 < |\eta| < 1.52$). Jets are reconstructed using the anti- k_t algorithm using all final-state stable particles as input, except those that are part of a dressed-lepton object. Jets are required to have $p_T > 25 \text{ GeV}$ and $|y| < 4.4$. Using these jets and leptons, events are then selected in a VBF topology using requirements identical to those imposed at detector-level. The VBF topology, SR and the three CRs are defined in Table 5.5.

Dressed muons	$p_T > 25 \text{ GeV}$ and $ \eta < 2.4$
Dressed electrons	$p_T > 25 \text{ GeV}$ and $ \eta < 2.37$ (excluding $1.37 < \eta < 1.52$)
Jets	anti- k_t , $p_T > 25 \text{ GeV}$ and $ y < 4.4$
VBF topology	$N_\ell = 2$ (same flavour, opposite charge), $m_{\ell\ell} \in (81, 101) \text{ GeV}$ $\Delta R_{\min}(\ell_1, j) > 0.4$, $\Delta R_{\min}(\ell_2, j) > 0.4$ $N_{\text{jets}} \geq 2$, $p_T^{j1} > 85 \text{ GeV}$, $p_T^{j2} > 80 \text{ GeV}$ $m_{jj} > 1 \text{ TeV}$, $ \Delta y_{jj} > 2.0$, $p_T^{\ell\ell} > 20 \text{ GeV}$, $p_T^{\text{bal}} < 0.15$, $\xi_Z < 1.0$
CRA	$\text{VBF topology} \oplus N_{\text{jets}}^{\text{gap}} \geq 1$ and $\xi_Z < 0.5$
CRb	$\text{VBF topology} \oplus N_{\text{jets}}^{\text{gap}} \geq 1$ and $\xi_Z > 0.5$
CRc	$\text{VBF topology} \oplus N_{\text{jets}}^{\text{gap}} = 0$ and $\xi_Z > 0.5$
SR	$\text{VBF topology} \oplus N_{\text{jets}}^{\text{gap}} = 0$ and $\xi_Z < 0.5$

TABLE 5.5: Fiducial definition of the measurement. $\Delta R_{\min}(\ell_1, j)$ denotes the minimum ΔR distance between the highest transverse-momentum lepton (ℓ_1) and any of the jets in the event. $\Delta R_{\min}(\ell_2, j)$ is defined analogously.

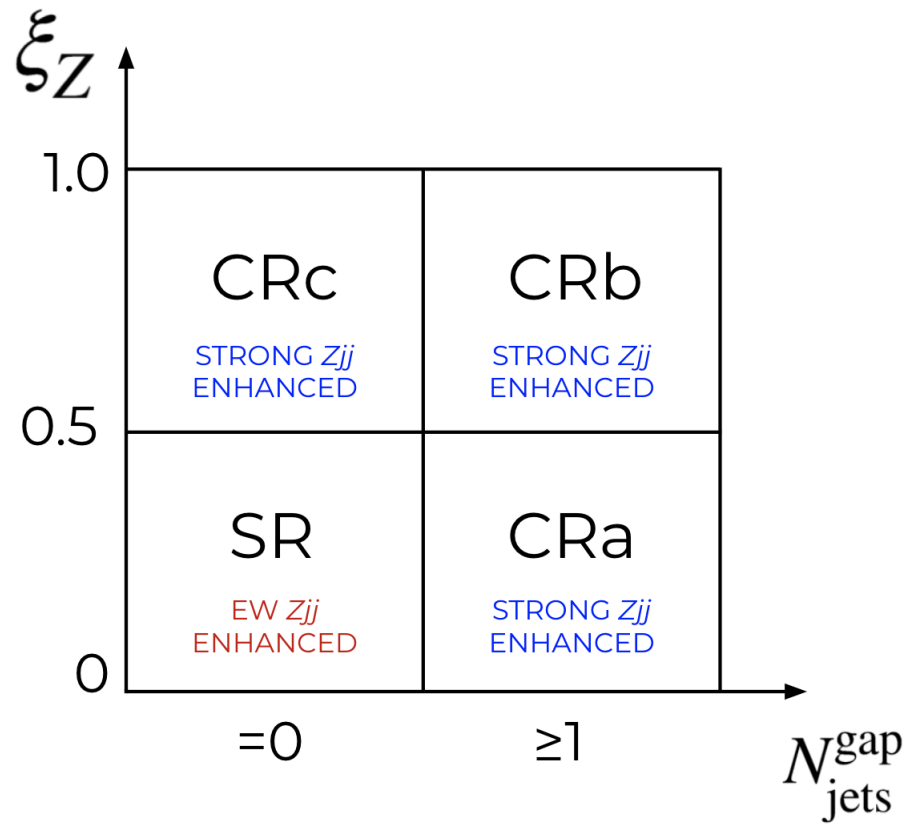


FIGURE 5.14: Diagram depicting how the control and signal regions are defined for the extraction of the EW- Z_{jj} component.

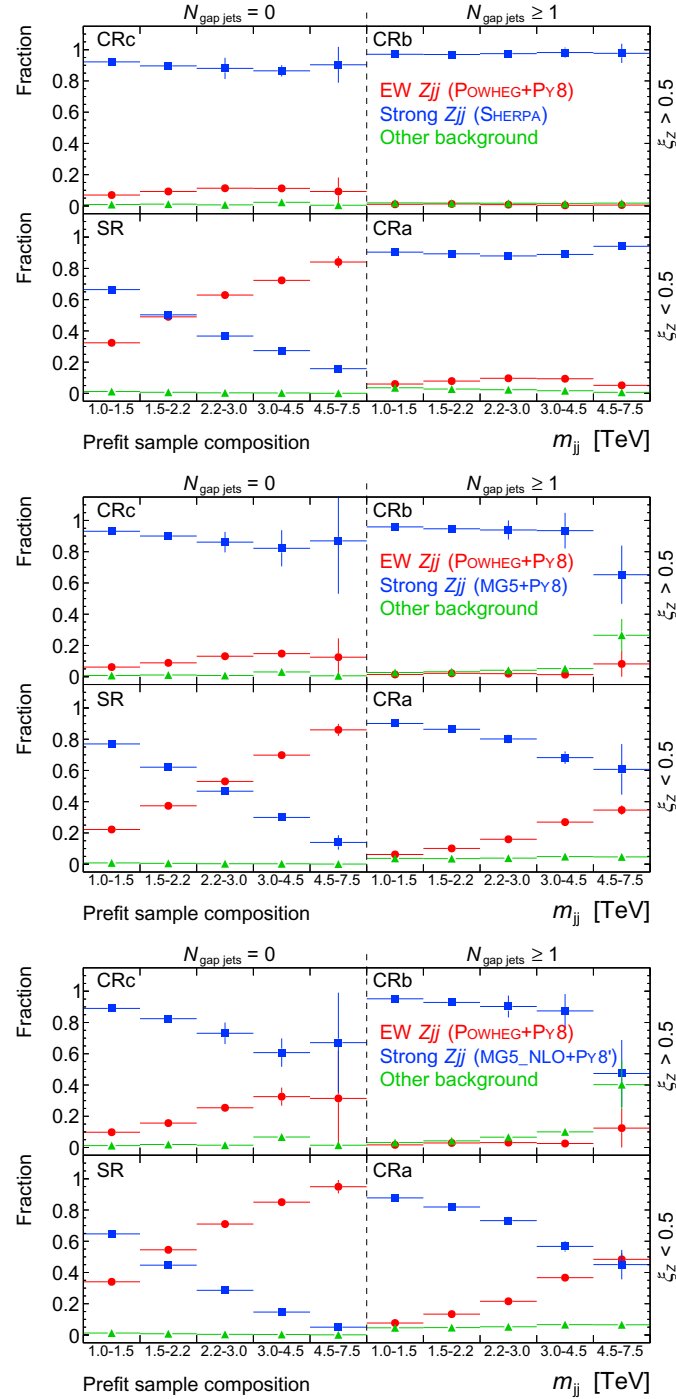


FIGURE 5.15: Predicted sample composition in m_{jj} bins of the signal region and control regions. The SM-predicted EW- Zjj fraction is shown as red circle, the fraction of strong- Zjj in blue squares, and the small contribution from other sources is shown as green triangles. The upper figure use predictions of strong- Zjj from SHERPA, the middle from MG5+PY8 and the bottom from MG5_NLO+PY8'. The latter generator exhibits the best agreement with the data.

5.4 Binning of the measured observables

The binning of the observables used for the differential measurements is another important consideration. The binning used is presented in Table 5.6, the first four bins ($m_{jj} < 1000$ GeV) of the m_{jj} variable are only studied for the inclusive Zjj measurement as a function of m_{jj} . For all other measurements a requirement of $m_{jj} > 1000$ GeV is applied. The final bin in m_{jj} and $|\Delta y_{jj}|$ have low statistics in CRb and CRc, so for some purposes the two last bins of m_{jj} and $|\Delta y_{jj}|$ are merged.

Observable	Bin edges											
m_{jj} / GeV	250	350	500	750	1000	1500	2250	3000	4500	7500		
$ \Delta y_{jj} $	2.0	2.6	3.1	3.6	3.9	4.2	4.6	5.0	6.0	8.0		
$p_T^{\ell\ell} / \text{GeV}$	20	30	45	70	100	140	200	275	400	550	1050	
$\Delta\phi_{jj}^{\text{signed}}$	$-\pi$	$-\frac{15\pi}{16}$	$-\frac{7\pi}{8}$	$-\frac{3\pi}{4}$	$-\frac{\pi}{2}$	$-\frac{\pi}{4}$	0	$\frac{\pi}{4}$	$\frac{\pi}{2}$	$\frac{3\pi}{4}$	$\frac{7\pi}{8}$	$\frac{15\pi}{16}$ π

TABLE 5.6: Binning used for the Zjj differential cross section measurements.

The binnings are chosen to satisfy the following criteria:

- Bins should have similar statistical precision across all measurement variables and channels. This criteria is relaxed at the extremes of some variables.
- The absolute detector resolution of a given measurement should be smaller than each bin of the distribution. Assuming Gaussian response the width of the bin should be greater than the 2σ deviation of the response distribution, where the detector response is the ratio of the quantity at detector-level to particle-level. This is to ensure that an event placed in a given bin is measured with a resolution that is small compared to the bin width.
- In steeply falling distributions ($m_{jj}, p_T^{\ell\ell}$), bins should have roughly even logarithmic spacing.

- Bin edge values should be chosen to ease comparison and presentation of the results.

For a variable such as $\Delta\phi_{jj}^{\text{signed}}$, the range is defined by geometry to be $[-\pi, \pi]$. The other 3 variables (m_{jj} , $p_T^{\ell\ell}$ and $|\Delta y_{jj}|$) do not have a clear upper bound so one is determined by observing where the MC prediction falls to a certain number of predicted events. In particular, the upper bound is placed at a round value nearby the value k where the integral from $[k, \infty] = 1$ (*i.e.* one predicted event for variable $> k$). The lower bound of this first bin is placed nearby k where the integral over the range $[k, \infty] = 30$, this creates a bin with sufficient expected events that the statistical uncertainty will not be overwhelming but that the bin is still narrow enough to be of interest. These expected precisions are determined using the SHERPA strong- Z_{jj} prediction reweighted to the shape and yield of the data in the control regions.

The distribution of $\Delta\phi_{jj}^{\text{signed}}$ reflects how the leading jets tend to be produced back to back, so the two bins around $\Delta\phi_{jj}^{\text{signed}} = 0$ will have the lowest statistical precision. It was determined that the narrowest of these bins should be $\pi/4$ to have sufficient expected statistics, then the bins are made narrower in even fractions of π as the value of $\Delta\phi_{jj}^{\text{signed}}$ approaches $\pm\pi$ to keep the statistical precision relatively constant. After the highest bin of m_{jj} , $p_T^{\ell\ell}$ and $|\Delta y_{jj}|$ was determined, another 7–9 bins are spaced out logarithmically and then slight adjustments are made to achieve a flat statistical precision.

Chapter 6

Measurements

This chapter describes in detail the technique used to perform the first measurements of EW- Zjj differential cross sections. These measurements are important for several reasons. First, we can study the distributions to test the SM predictions and see where they behave as expected. The measured spectra can be compared to the predictions from different MC generators and see which of these best describes the data. We can also go a step further and look for signs of new physics and compare the measurements to beyond the Standard Model predictions.

The analysis techniques employed in the EW- Zjj signal extraction are relatively complex when compared to previous analyses of the same process [5, 7]. The measurements are done using two main steps. First there is a binned maximum likelihood fit to derive a data driven correction, which is used to correct the strong- Zjj component. With the strong component corrected, it is then possible to extract the EW- Zjj cross section from the data. The second main step is to take this detector-level EW- Zjj cross section result and unfold the effects of the detector to obtain a particle-level result.

The data driven correction and EW signal extraction procedure is complex due

to the large strong- Zjj background that is poorly modelled by the MC simulation in the VBF topology phase space as shown in Section 5.2. The final method presented here is the result of years of work by myself and the analysis team, it went through many iterations and optimizations.

This chapter begins, in Section 6.1, by demonstrating the signal extraction method with a practical example. The method is then formally defined in Section 6.2. Then, in Section 6.2.1, further details about the choices made when designing the model are explained. The method is then validated via tests performed on simulated (Asimov) data in Section 6.3 and the various sources of systematic uncertainty are described in Section 6.4. The validated method is used to extract the EW- Zjj cross sections from the data, which are presented in Section 6.5. Then, in Section 6.6, several variations to the statistical method are introduced to study the impact they have on the measured result.

The measured EW- Zjj event yields at detector-level are then *unfolded*, meaning to correct for detector inefficiency and resolution in order to obtain a particle-level result. This correction procedure is described in Section 6.7 and the unfolded, particle-level results are presented in Section 6.8. These particle-level results are then compared to the EW- Zjj predictions from simulation, the first main goal of the analysis.

6.1 EW signal extraction method

To illustrate the strong- Zjj modelling problem, and how the fit method corrects it, an example fit to the ATLAS data in the m_{jj} observable will be presented. It

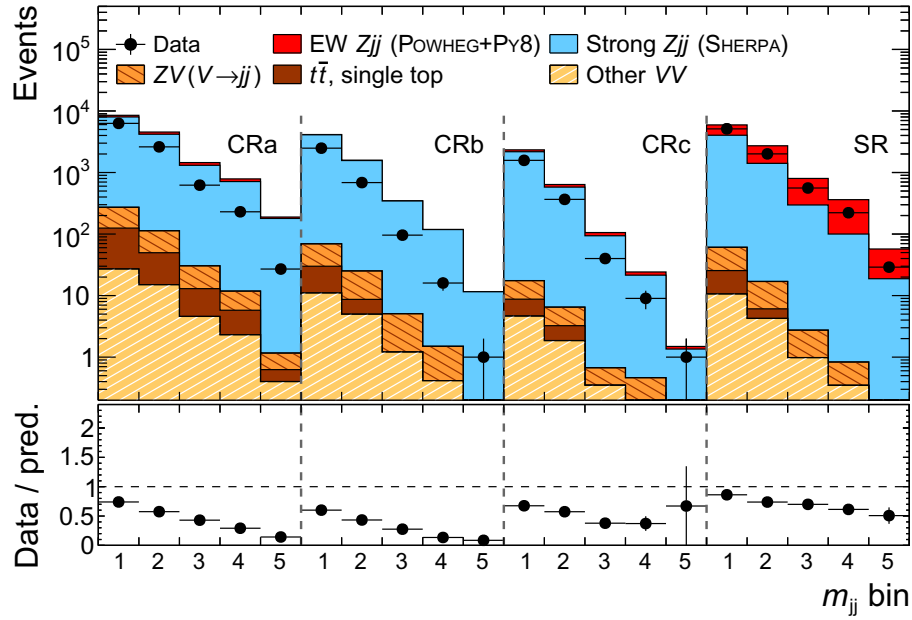


FIGURE 6.1: Pre-fit data to MC comparison in each of the four regions as a function of m_{jj} . The SHERPA strong- Zjj component over predicts the data in all bins.

is important to note that this example is meant to show the method in practice; however it is not representative of how the method was developed. In development, steps were taken to avoid biasing the results to the data. Preliminary tests and validations were performed with simulated Asimov data and when testing with real data the signal region was blinded.

To begin the example, the pre-fit data to MC comparison is shown in Figure 6.1. The SHERPA prediction is used for the strong- Zjj component, and as can be seen in the figure the MC simulation over predicts the data in all bins. The data to MC discrepancy increases as m_{jj} increases.

The description of this section is meant to introduce the method conceptually by breaking the fit down into three steps. The corrections obtained from these steps

will be to cumulatively applied to the MC predictions in Figure 6.1. In reality, these three steps occur simultaneously in a global fit, which will be precisely described in Section 6.2.

The first step in correcting the strong- Zjj in the SR is determined (primarily) from the data to MC ratio in a control region. The EW- Zjj component is suppressed in the control regions, thus the strong- Zjj process is isolated and it can be better matched to the data. It was found that CRa is best suited for this purpose, as will be shown in Section 6.2.1. There is one parameter, $b_{L,i}$, per bin i of the measurement. This parameter scales the strong- Zjj component in the SR and in CRa. The label “L” denotes the Low Z centrality criteria applied in the two regions ($\xi_Z < 0.5$).

This leaves two other control regions, CRb and CRc, whose purpose in the fit will be explained in the next step, but for now they are used to derive another data driven constraint. Just as CRa is primarily used to derive the $b_{L,i}$ constraints applied to CRa and to the SR, another set of constraints $b_{H,i}$ is derived primarily in CRb and applied to CRc. Similarly, “H” denotes the High Z centrality criteria applied in these two regions ($0.5 < \xi_Z < 1.0$). Figure 6.2 shows the data to simulation comparison with the fitted $b_{L,i}$ and $b_{H,i}$ factors applied to the strong- Zjj component in the appropriate regions. Note that the simulation shown in the following figures has been adjusted by the fitted factors and is no longer the default MC prediction. This figure shows that the data to MC in CRa and CRb are now in agreement, this is expected as these regions determine the corrections. The agreement in these regions is not perfect and this is expected since these $b_{L,i}$ and $b_{H,i}$ factors are the only the first corrections in an interconnected model. The agreement in the SR and in CRc has also been improved, so some residual correction is still necessary in all regions.

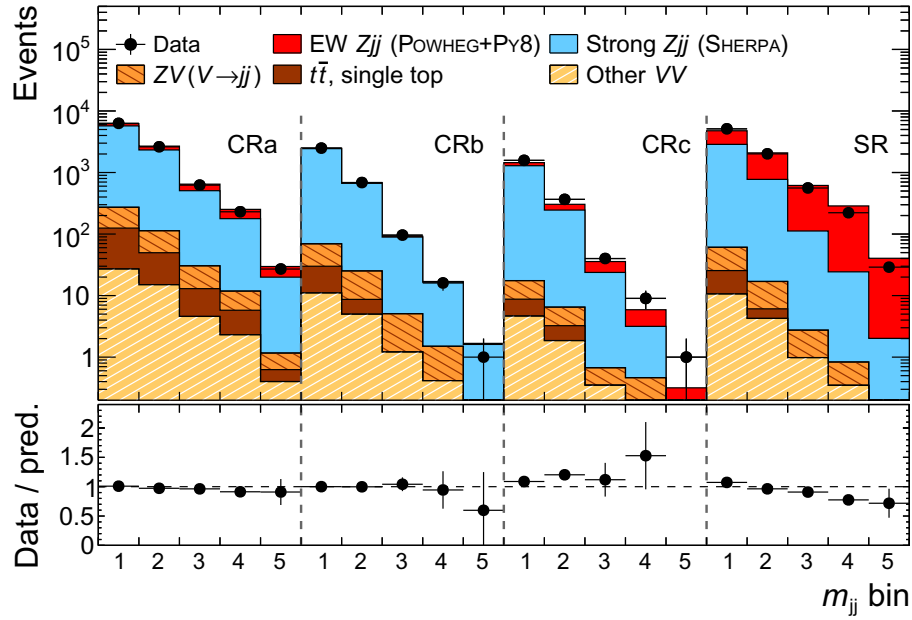


FIGURE 6.2: Data to simulation comparison after applying the $b_{L,i}$ and $b_{H,i}$ factors. These bin-by-bin factors account for the bulk of the correction applied to the strong- Zjj component.

The next step introduces this residual correction, using CRc to derive a function $f(x)$ (in this case $f(m_{jj})$) that will further correct the strong- Zjj component in both CRc and the SR. The form of this function was chosen to be a first-order polynomial, $f(x) = a|x| + b$. The absolute value is only of consequence for the $\Delta\phi_{jj}^{\text{signed}}$ observable, the other measurement variables are strictly positive. This residual correction serves to “link” the two pairs of regions such that the three control regions are all used to constrain the SR, this is illustrated in Figure 6.3.

The fitted $f(m_{jj})$ function is overlaid with the data to simulation ratio in CRc in Figure 6.4. The ratio is shown after applying the $b_{H,i}$ factors to the strong component in each bin, as in the lower panel of Figure 6.2. Figure 6.5 shows the data to simulation comparison after applying the $f(x)$ constraint to CRc and to the

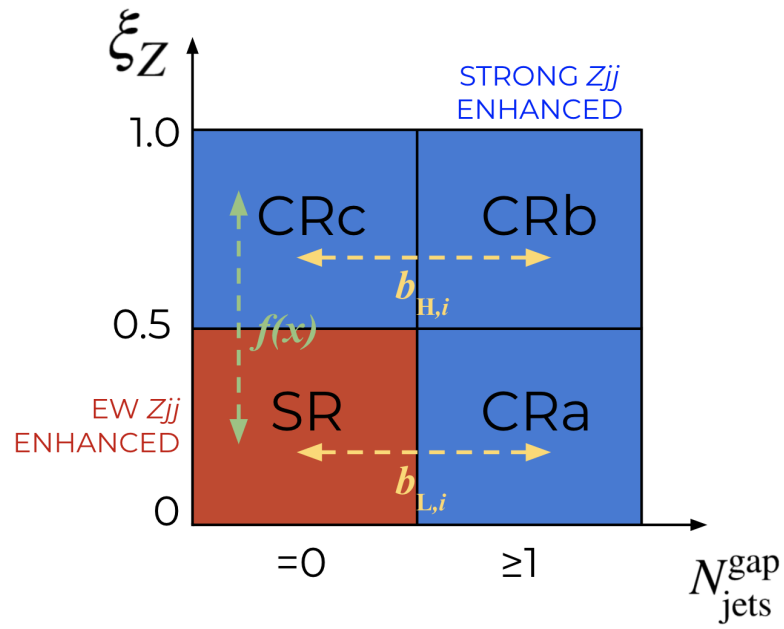


FIGURE 6.3: Diagram depicting how the $b_{L,i}$ / $b_{H,i}$ factors and the $f(x)$ function are applied to the signal and control regions to constrain the strong- Zjj component.

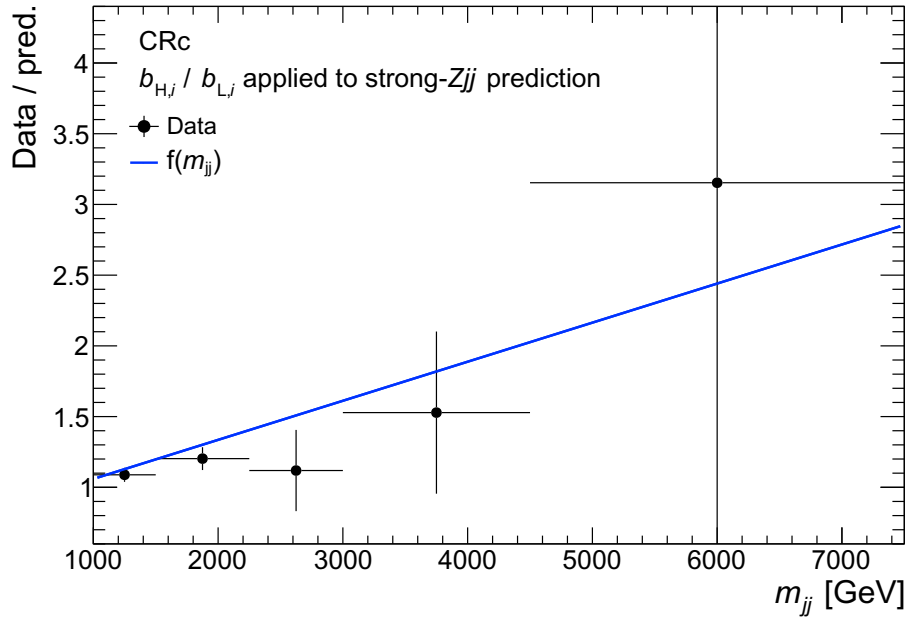


FIGURE 6.4: The fitted $f(m_{jj})$ function overlaid on the data to simulation ratio in CRc. This comparison is made after the bin-by-bin factors were applied to the strong- Zjj component.

SR. The value applied in each bin is determined by evaluating $f(x)$ at the central m_{jj} value of the bin. At the stage the strong- Zjj component has been corrected and there is good agreement in all control region bins. However there is still a remaining offset in the SR that will be corrected in the final step.

Since the strong component is now corrected, the final step is to extract the EW- Zjj signal defined in each bin of the signal region as data minus the strong- Zjj and other background components. The final post-fit data to simulation comparison is shown in Figure 6.6. The EW signal is measured in each bin with an independent free parameter that scales the EW component to match the data.

To recap, the figures shown in this section are meant as a conceptual demonstration of the method. In reality, all three steps of the fit happen simultaneously

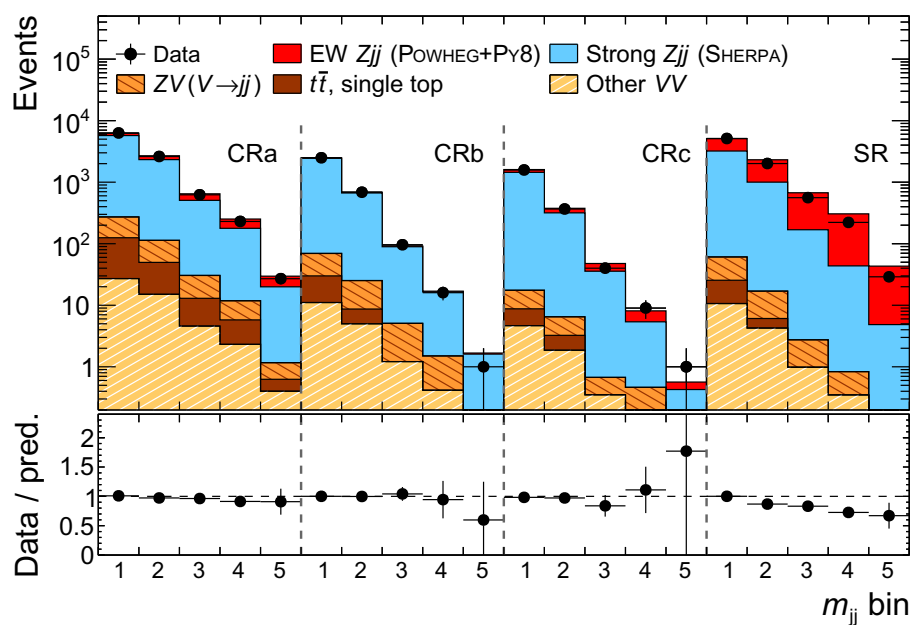


FIGURE 6.5: Data to simulation comparison after applying the $f(x)$ constraint. At this stage the strong- Zjj component has been corrected and any residual discrepancy will be corrected by adjusting the EW- Zjj component.

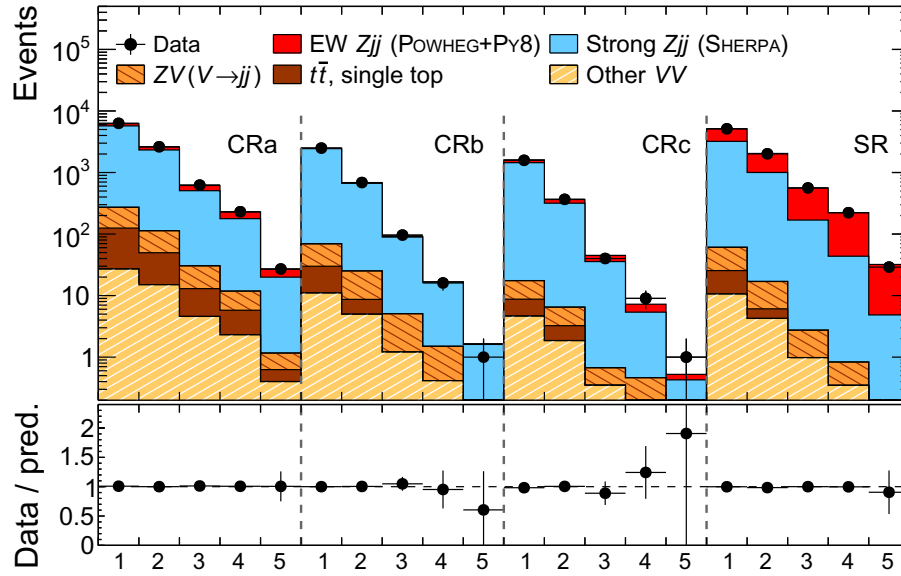


FIGURE 6.6: Post-fit data to simulation comparison after applying the μ_i parameters to the EW component. The data and simulation now agree to within the data statistical uncertainty in all bins.

in all regions, so constraints are not technically derived in one region and applied to another as described. However, it will be shown in the following sections that the fit model has been designed in such a way that its general behaviour matches what has been described here conceptually. In fact, the $b_{H,i}$, $b_{L,i}$ and $f(m_{jj})$ factors applied in the figures were obtained simultaneously with the final likelihood fit method. Applying them one at a time to the MC distributions is to help illustrate the role each constraint plays in the overall method.

6.2 Statistical method definition

The four regions, introduced in Section 5.3, enter a binned maximum likelihood fit to constrain the strong- Zjj prediction in the signal region via constraints defined

across multiple regions. The fit is performed using the TROOFIT package [106]. The fitted number of events in each region and in each bin of a distribution is given by

$$\nu_{ri} = \mu_i \nu_{ri}^{\text{EW,MC}} + \nu_{ri}^{\text{strong}} + \nu_{ri}^{\text{other,MC}}, \quad (6.1)$$

where μ_i is the EW- Zjj signal strength of bin i , a correction factor applied to the signal MC. r is an index corresponding to the region $r \in \{\text{CRa,CRb,CRc,SR}\}$. ν_{ri}^{strong} is the corrected strong- Zjj component. The majority of the parameters fit in this method are designed to correct the strong- Zjj MC prediction. $\nu_{ri}^{\text{EW,MC}}$ and $\nu_{ri}^{\text{other,MC}}$ are the MC predictions of EW- Zjj and contributions from “other” processes (diboson, $t\bar{t}$ and single top), respectively. The EW yield to be measured in the signal region is defined as

$$\nu_{\text{SR},i}^{\text{EW}} = \mu_i \nu_{\text{SR},i}^{\text{EW,MC}}. \quad (6.2)$$

The signal strength parameter μ_i is applied in the fit to not only the EW- Zjj in the signal region but also to the small EW component in the control regions. It is defined this way in order to correlate the measured yield in the signal region with the small EW- Zjj yield in the control regions.

In order to extract the EW signal, the strong- Zjj estimation is constrained using the different EW-suppressed control regions according to

$$\begin{aligned}
 \nu_{\text{CRa},i}^{\text{strong}} &= b_{\text{L},i} \nu_{\text{CRa},i}^{\text{strong,MC}}, \\
 \nu_{\text{CRb},i}^{\text{strong}} &= b_{\text{H},i} \nu_{\text{CRb},i}^{\text{strong,MC}}, \\
 \nu_{\text{CRc},i}^{\text{strong}} &= b_{\text{H},i} f(x_i) \nu_{\text{CRc},i}^{\text{strong,MC}}, \\
 \nu_{\text{SR},i}^{\text{strong}} &= b_{\text{L},i} f(x_i) \nu_{\text{SR},i}^{\text{strong,MC}}.
 \end{aligned} \tag{6.3}$$

Here, the $b_{\text{L},i}$ and $b_{\text{H},i}$ are sets of bin-dependent factors that apply to the low ($\xi_Z < 0.5$) and high ($0.5 < \xi_Z < 1.0$) Z-centrality regions, respectively. These factors are primarily constrained in CRa and CRb, where they adjust the predicted simulated strong- Zjj event yields and bring the total predicted yield in better agreement with data. The $f(x_i)$ is the value of a two-parameter function of the observable x being measured evaluated at the centre of bin i (x_i). This function provides a residual correction to the constrained strong- Zjj yield to account for the extrapolation from CRa ($N_{\text{jets}}^{\text{gap}} \geq 1$) and the SR ($N_{\text{jets}}^{\text{gap}} = 0$) and is primarily constrained by CRc. The function form was chosen to be a first-order polynomial, $f(x) = a|x| + b$. Figure 6.3 summarizes how the strong- Zjj constraints are applied to the four regions.

The log likelihood is defined according to

$$\ln L(\mathbf{N}^{\text{data}}; \boldsymbol{\theta}) = - \sum_{r,i} \nu_{ri}(\boldsymbol{\theta}) + \sum_{r,i} N_{ri}^{\text{data}} \ln \nu_{ri}(\boldsymbol{\theta}) - \sum_s \frac{\theta_s^2}{2}, \tag{6.4}$$

where again r is an index corresponding to the region, i is the bin of the kinematic observable, N_{ri}^{data} is the observed event yield and $\nu_{ri}(\boldsymbol{\theta})$ is the prediction that is

dependent on the sources of systematic uncertainty that are each constrained by *nuisance parameters* (NPs) $\theta = (\theta_1, \dots, \theta_s)$. The NPs will modify the ν_{ri} quantities in the fit according to the fractional size of uncertainty source j in bin i of the observable, δ_{ij} . These δ_{ij} will differ for the strong and EW component of the total prediction as the uncertainty sources effect each process differently. The sign of the δ_{ij} terms determine the correlation between bins and regions as well as between the different processes. Combining all NPs θ_j gives a modification for the strong- Zjj MC prediction of the form

$$\nu_{ri}^{\text{strong}} \cdot \prod_j \left(1 + \theta_j \delta_{ij}^{\text{strong}}\right). \quad (6.5)$$

Here the $\delta_{ij}^{\text{strong}}$ are particular to the strong process. The EW- Zjj and other background MC predictions will be modified analogously by their own set of NPs corresponding to the relevant uncertainty sources for each process.

The fit minimizes the likelihood function to determine the best fit values of the model parameters such as the parameters of interest, the EW signal strengths μ_i . The extracted EW- Zjj yield is then $\mu_i \nu_{ri}^{\text{EW,MC}}$.

The $b_{L,i}$ and $b_{H,i}$ factors are shared between pairs of bins with the same absolute $\Delta\phi_{jj}^{\text{signed}}$, thus there are 6 such parameters to constrain the 12 $\Delta\phi_{jj}^{\text{signed}}$ bins. The strong- Zjj process has been studied extensively [35] and is symmetric as a function of $\Delta\phi_{jj}^{\text{signed}}$. Conversely, there are 12 μ_i EW signal strength parameters for $\Delta\phi_{jj}^{\text{signed}}$, one per measurement bin as for the other three observables, as these parameters represent the measured EW- Zjj result.

The free parameters in the binned maximum likelihood fit are therefore the

signal strengths μ_i , the two parameters of the function $f(x)$, and the $b_{L,i}$ and $b_{H,i}$ corrections to the strong- Zjj process. In total, this amounts to $3 N_{\text{bins}} + 2$ parameters that are constrained using $4 N_{\text{bins}}$ measurements in data. There is however an exception, the expected and observed yield in the final bin of m_{jj} and $|\Delta y_{jj}|$ in CRb and CRc is very low ($\mathcal{O}(1)$ event) so in these regions for these variables the bin dependent factor is merged between the last two bins.

The value of the parameters that constrain the strong- Zjj component (the $f(x)$ parameters, the $b_{L,i}$ and $b_{H,i}$) essentially determine the overall result. The remaining μ_i will control the normalization of the EW- Zjj component and, since there is one μ_i per bin the fit will prefer the value that allows for good data to MC agreement in the SR where the vast majority of EW- Zjj events reside. However the μ_i are applied to all regions, thus the small EW component in the CRs will have an impact on the μ_i parameters but to a much lesser extent than the SR.

As described in Section 4.1.2 there are three strong- Zjj samples, SHERPA and MG5+PY8 are both available at reconstructed level, but the MG5_NLO+PY8 is only available at particle-level. By applying a truth level reweighting to the nominal MG5+PY8 strong- Zjj simulation. This reweighting assigns a weight to each event such that the particle-level distributions of MG5+PY8 agree with the more precise MG5_NLO+PY8. This adjustment to the event weights also affect the reconstructed level distributions, which allows this sample to be included as a third strong- Zjj prediction for the fit. This reweighted sample is referred to as MG5_NLO+PY8'. The reweighting procedure is detailed in Appendix A.

There is no *a priori* reason to prefer any strong- Zjj event generator over another. Thus, the EW- Zjj component is extracted three times, once using the SHERPA strong-

Zjj prediction, once using the MG5+PY8 strong- Zjj and once using the MG5_NLO+PY8 reweighted strong- Zjj prediction. The final electroweak signal yield in each bin of the differential distribution is taken to be the midpoint of the envelope of yields obtained using the three different strong- Zjj event generators and the envelope itself defines the systematic uncertainty associated with the choice of event generator. This strong- Zjj generator choice systematic is the dominant source of systematic uncertainty in most bins of the measurement.

The constraints on the strong- Zjj simulation in Equation (6.1) are evaluated independently for each of the measured differential distributions ($m_{jj}, |\Delta y_{jj}|, p_T^{\ell\ell}$ and $\Delta\phi_{jj}^{\text{signed}}$). This causes slightly different total EW- Zjj and strong- Zjj event yields when summed across each differential spectrum. To ensure consistency between the distributions, an additional constraint is applied in the likelihood to ensure the same integrated strong- Zjj yield is obtained for each distribution, *i.e.*

$$\sum_i \nu_{\text{SR},i}^{\text{strong}} = \hat{\nu}_{\text{SR},m_{jj}}^{\text{strong}} \quad (6.6)$$

where $\hat{\nu}_{\text{SR},m_{jj}}^{\text{strong}}$ is the event yield obtained by integrating the constrained strong- Zjj template for the m_{jj} distribution in the SR. The choice of m_{jj} is motivated by the fact that the EW and strong Zjj processes have very different shapes as a function of m_{jj} . As a consequence, the m_{jj} -based measurement of the integrated strong- Zjj event yield is the most precise.

6.2.1 Fit method motivation

This section is to motivate the choice of the regions used to constrain the model parameters described in Section 6.1, the main choice being which control region (CRa, CRb or CRc) will share the set of bin-dependent factors (b_i) with the SR. CRb is an unnatural choice because in CRb both the $N_{\text{jets}}^{\text{gap}}$ and ξ_Z selection criteria are inverted with respect to the SR, so a bin-dependent factor constraining CRb and the SR would be said to *cross the boundary* of both $N_{\text{jets}}^{\text{gap}}$ and ξ_Z . Alternatively, using CRa or CRc for the bin-dependent factors crosses only one such boundary. In short, the selection criteria of CRa and CRc is closer to that of the SR than CRb is to the SR, so these regions are better suited to constrain the bin-dependent factors.

These leaves two choices, CRa or CRc, to share b_i factors with the SR. Once this choice is established, for example in the nominal method of Section 6.1 where CRa and the SR share the $b_{L,i}$ factors, the $f(x)$ function should then be applied in the CRc region and the SR to provide the residual correction. This is motivated by the notion that the selection relating CRa to the SR is similar to the selection relating CRb to CRc, the only difference between the two pairs of regions is the inversion of the $N_{\text{jets}}^{\text{gap}}$ criteria. Likewise if CRc is instead chosen to share b_i with the SR, then the $f(x)$ function should be applied to CRa and the SR.

With the four regions defined as they are, there are two reasonable ways to set up a model of the type described, either the nominal method with the b_i shared between CRa and the SR (CRa method) or the alternative method with b_i shared between CRc and the SR (CRc method). In each case the $f(x)$ function is applied to the SR and whichever of CRa or CRc that does not share the b_i with the SR. The CRa and CRc methods initially seem like they might perform differently in constraining

the strong- Zjj component of the SR, however in practice it will be shown that the two methods are essentially equivalent in this respect due to the symmetry of the defined regions and model parameters in each case.

To show this it is useful to introduce a quantity $X(r, i)$, where this r and i are the same region and bin indices defined above. Assuming the nominal CRa method and starting from Equation (6.3) for CRa, the fitted strong- Zjj is $\nu_{\text{CRa},i}^{\text{strong}} = b_{\text{L},i} \nu_{\text{CRa},i}^{\text{strong,MC}}$ and this is fitted to the strong- Zjj component of the data, *i.e.* postfit the relationship will be

$$\nu_{\text{CRa},i}^{\text{strong}} = b_{\text{L},i} \nu_{\text{CRa},i}^{\text{strong,MC}} \approx N_{\text{CRa},i}^{\text{data}} - \nu_{\text{CRa},i}^{\text{EW,MC}} - \nu_{\text{CRa},i}^{\text{other,MC}}, \quad (6.7)$$

where the signal strength parameter μ_i has been set to 1, assuming the SM prediction. This assumption is only made here to simplify this conceptual explanation, the signal extraction method does not assume the SM. This approximate relation motivates defining the quantity $X(r)$ as

$$X(r, i) = \frac{N_{\text{r},i}^{\text{data}} - \nu_{\text{ri}}^{\text{EW,MC}} - \nu_{\text{ri}}^{\text{other,MC}}}{\nu_{\text{r},i}^{\text{strong,MC}}}, \quad (6.8)$$

and then $b_{\text{L},i} \approx X(\text{CRa}, i)$. The $b_{\text{L},i}$ will be constrained in the fit to the ratio $X(\text{CRa}, i)$ since the SR is much more flexible in the fit. This is because the SR bins have the $f(x)$ function applied and a large EW- Zjj component that will largely determine the μ_i parameters. In contrast the data to MC agreement in CRa will be almost entirely determined by the $b_{\text{L},i}$ values. The other set of bin-dependent factors, $b_{\text{H},i}$, will be constrained approximately to the quantity $X(\text{CRb}, i)$ since again the function $f(x)$ is applied to CRc giving an extra dimension of flexibility to fit to the strong- Zjj

component in CRc, thus $b_{H,i} \approx X(\text{CRb}, i)$.

Looking again at Equation (6.3), this time for CRc

$$\nu_{\text{CRc},i}^{\text{strong}} = b_{H,i} f(x_i) \nu_{\text{CRc},i}^{\text{strong,MC}} \approx N_{\text{CRc},i}^{\text{data}} - \nu_{\text{CRc},i}^{\text{EW,MC}} - \nu_{\text{CRc},i}^{\text{other,MC}}, \quad (6.9)$$

thus $X(\text{CRc}, i) \approx b_{H,i} f(x_i) \approx X(\text{CRb}, i) f(x_i)$. The residual correction function is fit to the ratio $X(\text{CRc}, i) / X(\text{CRb}, i)$, combining this with $b_{L,i} \approx X(\text{CRa}, i)$ the correction to the SR in the CRa method as

$$b_{L,i} f(x_i) \approx \frac{X(\text{CRa}, i) X(\text{CRc}, i)}{X(\text{CRb}, i)}. \quad (6.10)$$

This is useful because similar reasoning shows that the CRc method gives the same ratio of $X(r, i)$. In this alternative CRc method the b_i shared with the SR will be approximately equal to $X(\text{CRc}, i)$ and the residual correction function is fit to the ratio $X(\text{CRa}, i) / X(\text{CRb}, i)$, multiplying these two constraints gives the correction applied to the SR which is again the same ratio $X(\text{CRc}, i) X(\text{CRa}, i) / X(\text{CRb}, i)$. The following section will show that this is what the model does in practice in a fit to simulated Asimov data.

6.3 Method validation

The signal extraction method has been validated with two simulated (*Asimov*) datasets, created by summing the POWHEG+PY8 EW- Zjj prediction, the small backgrounds (diboson, $t\bar{t}$ and single top) and a prediction for the strong- Zjj component. The Asimov datasets enter as real data in the likelihood fit, however since these data

are created from predictions, the data counts will have non-integer values. The first such dataset, referred to as “AsimovFxFx”, uses the MG5_NLO+PY8 prediction for the strong component. This prediction is the reconstructed level leading order MG5+PY8 prediction with a particle-level reweighting applied. As can be seen in Appendix A, this reweighted prediction achieves agreement with the data so the AsimovFxFx dataset can be considered realistic. “FxFx” refers to the matching and merging scheme used in the MG5_NLO+PY8 simulation.

The second Asimov dataset, referred to as “AsimovSrew” (Asimov Sherpa reweighted) uses the SHERPA strong- Zjj prediction that has been reweighted by smooth functions to match the data in the control regions via the procedure described in Appendix A. This is necessary since the nominal SHERPA prediction is in poor agreement with the data in m_{jj} and $|\Delta y_{jj}|$ (see Figure 5.11). Figures 6.7, 6.8 and 6.9 compare the real data to AsimovFxFx and AsimovSrew in CRa, CRb and CRc, respectively. The Asimov datasets both agree with the data within uncertainty for most bins, the exception being AsimovFxFx for $p_T^{\ell\ell}$ in CRb and CRc. As such the result from AsimovSrew will be more realistic for $p_T^{\ell\ell}$.

From Figures 6.7–6.9 it can be seen that the observed yield in the last bin m_{jj} and $|\Delta y_{jj}|$ is very small ($\mathcal{O}(1)$ event) in CRb and CRc, therefore a bin dependent factor derived in the last bin of these regions will do a poor job of constraining the strong- Zjj yield. To account for this, the final b_i is shared between the last two bins of m_{jj} and $|\Delta y_{jj}|$ in CRb and CRc. If the model is set up in the CRa method described in section 6.2.1 then the SR can be constrained with one bin dependent factor per bin from CRa, however the CRc method would need to merge the constraint applied to the SR in the last two bins of these variables. For this reason the choice was made

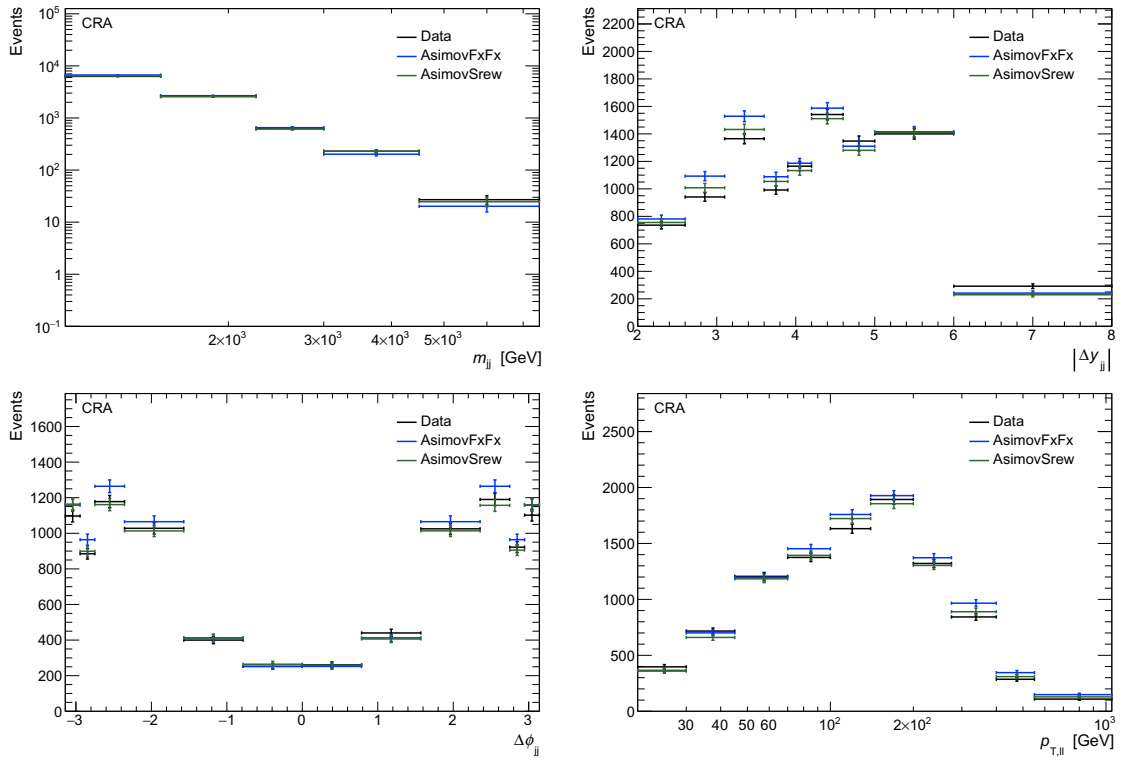


FIGURE 6.7: Comparison of Asimov datasets to real data for each measurement variable in CRA.

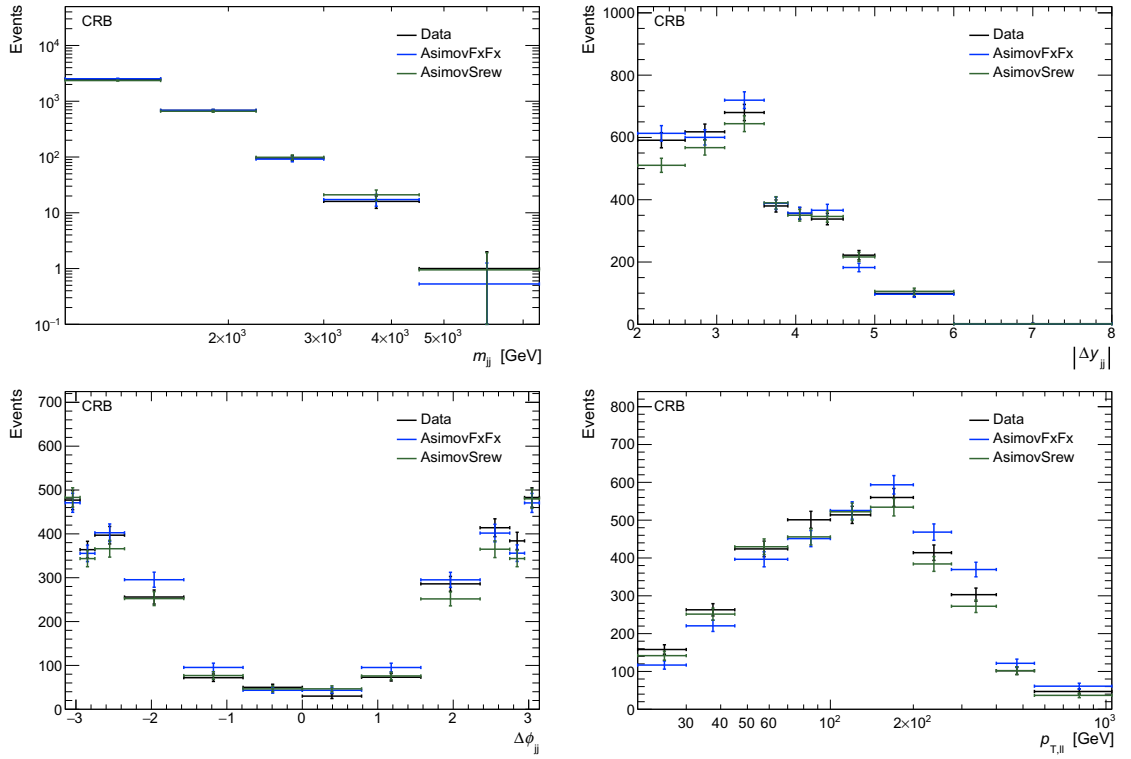


FIGURE 6.8: Comparison of Asimov datasets to real data for each measurement variable in CRb.

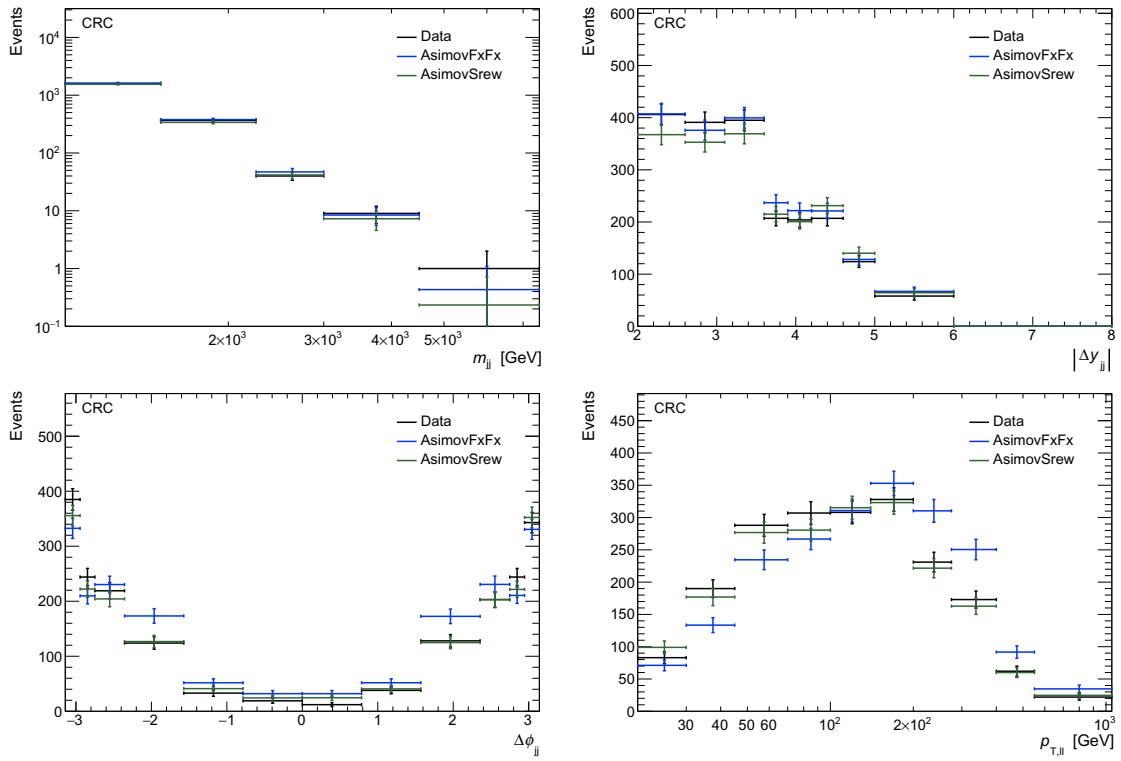


FIGURE 6.9: Comparison of Asimov datasets to real data for each measurement variable in CRc.

to use the CRa method, and it will be shown that in practice the two methods will produce a similar constraint in the SR.

The m_{jj} fit to AsimovSrew using the MG5_NLO+PY8 strong- Zjj template, with both the CRa and CRc method of implementing the model constraints is used to illustrate the performance of the model in practice. For the nominal CRa method, there are bin dependent factors $b_{L,i}$ and $b_{H,i}$ that apply to the low ($\xi_Z < 0.5$, SR and CRa) and high ($\xi_Z > 0.5$, CRb and CRc) Z-centrality regions, respectively. For the alternative CRc method there are bin dependent factors $b_{NGJ,i}$ and $b_{GJ,i}$ that apply to the “no gap jet” (SR and CRc) and the “gap jet” (CRa and CRb) regions, respectively. Recall that in CRa method only the $b_{H,i}$ need to have a merged final two bin constraint in m_{jj} and $|\Delta y_{jj}|$ since they are constrained in CRb and CRc, however in CRc method both $b_{NGJ,i}$ and $b_{GJ,i}$ need to have a merged final two bin constraint since they are constrained in CRc and CRb, respectively.

In Table 6.1 the computed $X(r)$ ratios are compared to the fitted constraints in both the CRa and CRc method, the values are similar as expected from the description in Section 6.2.1. The CRa and CRc methods produce similar EW signal strengths as expected, thus the methods constrain the strong- Zjj component of the SR in an equivalent way, so the choice to use CRa for the bin dependent factors shared with the SR as the nominal method is made because CRa has more statistics than CRc, particularly in the last bin of m_{jj} and $|\Delta y_{jj}|$.

With the CRa method established as the nominal method, the Asimov data can be used to test the bias of the EW extraction method by comparing the measured and injected EW- Zjj component. As explained in Section 6.1, the EW- Zjj component is extracted three times, once using the SHERPA strong- Zjj prediction, once using

	m_{jj} [TeV]	1 - 1.5	1.5 - 2.25	2.25 - 3.0	3.0 - 4.5	4.5 - 7.5
CRa method	$b_{H,i}$	0.93	0.96	1.09	1.28	1.28
	$X(\text{CRb})$	0.93	0.96	1.10	1.24	2.73
	$b_{L,i}$	0.96	0.96	0.94	1.27	1.49
	$X(\text{CRa})$	0.95	0.94	0.93	1.25	1.49
	$f(x)$	1.04	0.93	0.79	0.60	0.21
	$X(\text{CRc}) / X(\text{CRb})$	1.00	0.93	0.90	0.62	0.21
	$f(x) \cdot b_{L,i}$	1.00	0.89	0.75	0.76	0.32
CRc method	$X(\text{CRa}) X(\text{CRc}) / X(\text{CRb})$	0.97	0.85	0.71	0.79	0.17
	μ^{EW}	0.87	0.90	0.94	0.97	1.00
	$b_{GJ,i}$	0.93	0.96	1.00	1.38	1.38
	$X(\text{CRb})$	0.93	0.96	1.10	1.24	2.73
	$b_{\text{NGJ},i}$	0.96	0.90	0.91	0.77	0.77
	$X(\text{CRc})$	0.95	0.87	0.84	0.78	0.31
	$f(x)$	1.03	1.00	0.97	0.92	0.83
	$X(\text{CRa}) / X(\text{CRb})$	1.02	0.98	0.85	1.01	0.55
	$f(x) \cdot b_{\text{NGJ},i}$	0.99	0.90	0.88	0.71	0.64
	$X(\text{CRa}) X(\text{CRc}) / X(\text{CRb})$	0.97	0.85	0.71	0.79	0.17
	μ^{EW}	0.89	0.89	0.89	0.98	1.01

TABLE 6.1: Table comparing $X(r)$ ratios computed from data and MC templates to fitted constraints for the m_{jj} fit to AsimovSrew using the MG5_NLO+PY8' strong- Zjj MC template. Recall that constraints derived in the last two bins are merged for CRb and CRc due to limited statistics in the final bin of these regions.

the MG5+PY8 strong- Zjj and once using the MG5_NLO+PY8' reweighted strong- Zjj prediction. The final electroweak signal yield in each bin of the differential distribution is taken to be the midpoint of the envelope of yields obtained using the three different strong- Zjj event generators and the envelope itself defines the systematic uncertainty associated with the choice of strong generator.

The Asimov data are used to check that the uncertainty assigned to the choice of strong generator covers the observed bias in the extracted EW- Zjj component. Figures 6.10 and 6.11 show the extracted EW- Zjj component from each of the three fits with the three strong generators relative to the injected POWHEG+PY8 EW- Zjj prediction. Also shown is the coverage of the strong generator choice systematic, in the majority of bins the systematic covers the injected value of EW- Zjj . The statistical and strong generator choice systematics shown here are the values obtained from the singular fit to the nominal Asimov dataset.

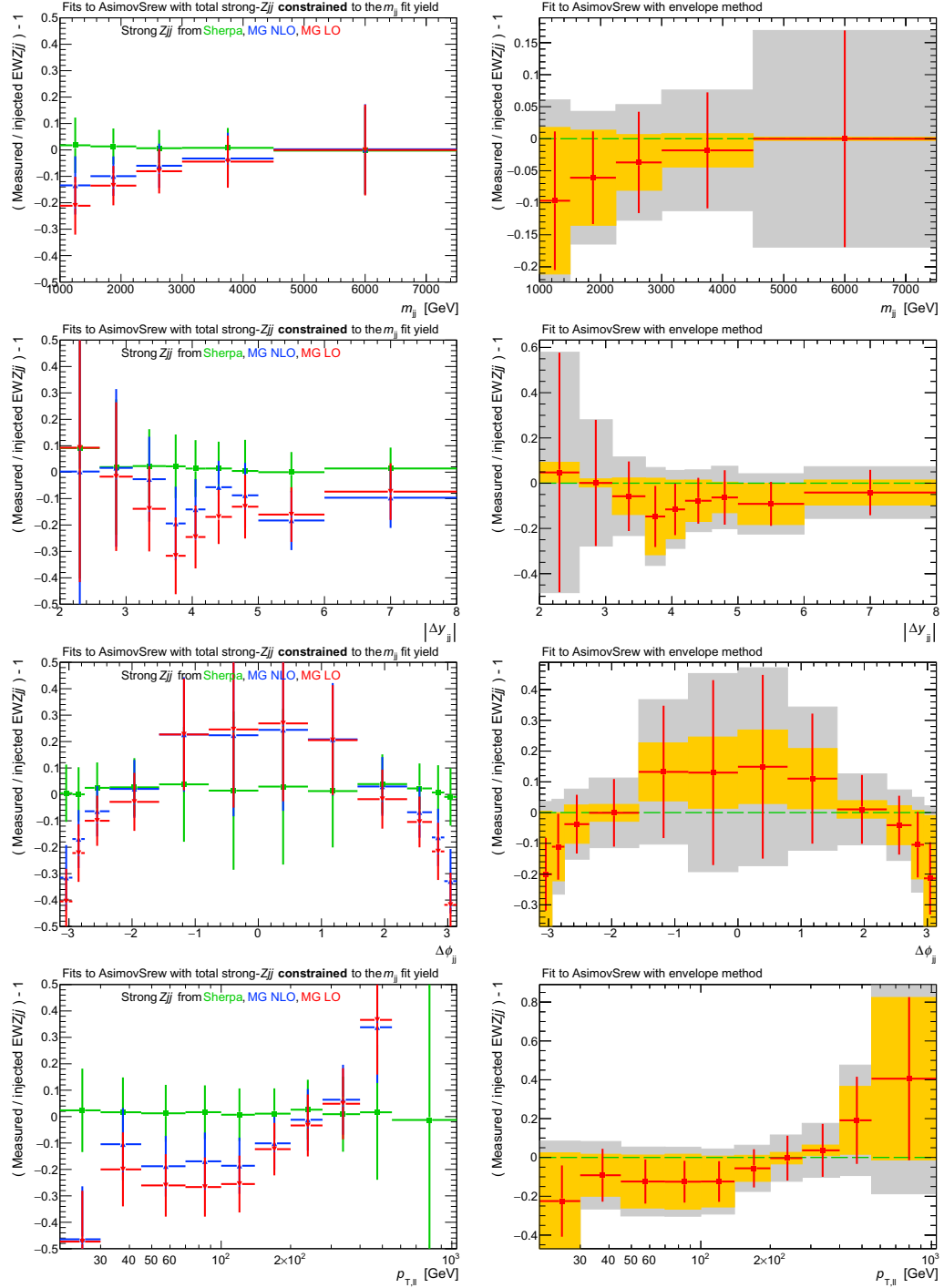


FIGURE 6.10: Measured EW- Zjj component for fits to AsimovSrew with each strong- Zjj generator relative to the injected POWHEG+PY8 EW- Zjj prediction (left) and coverage of the strong- Zjj generator choice systematic (orange) from the envelope of the three EW extraction fits (right). The statistical uncertainty from the likelihood fit is shown as a gray band, this is the numerical average of the statistical uncertainty of each of the three fits.

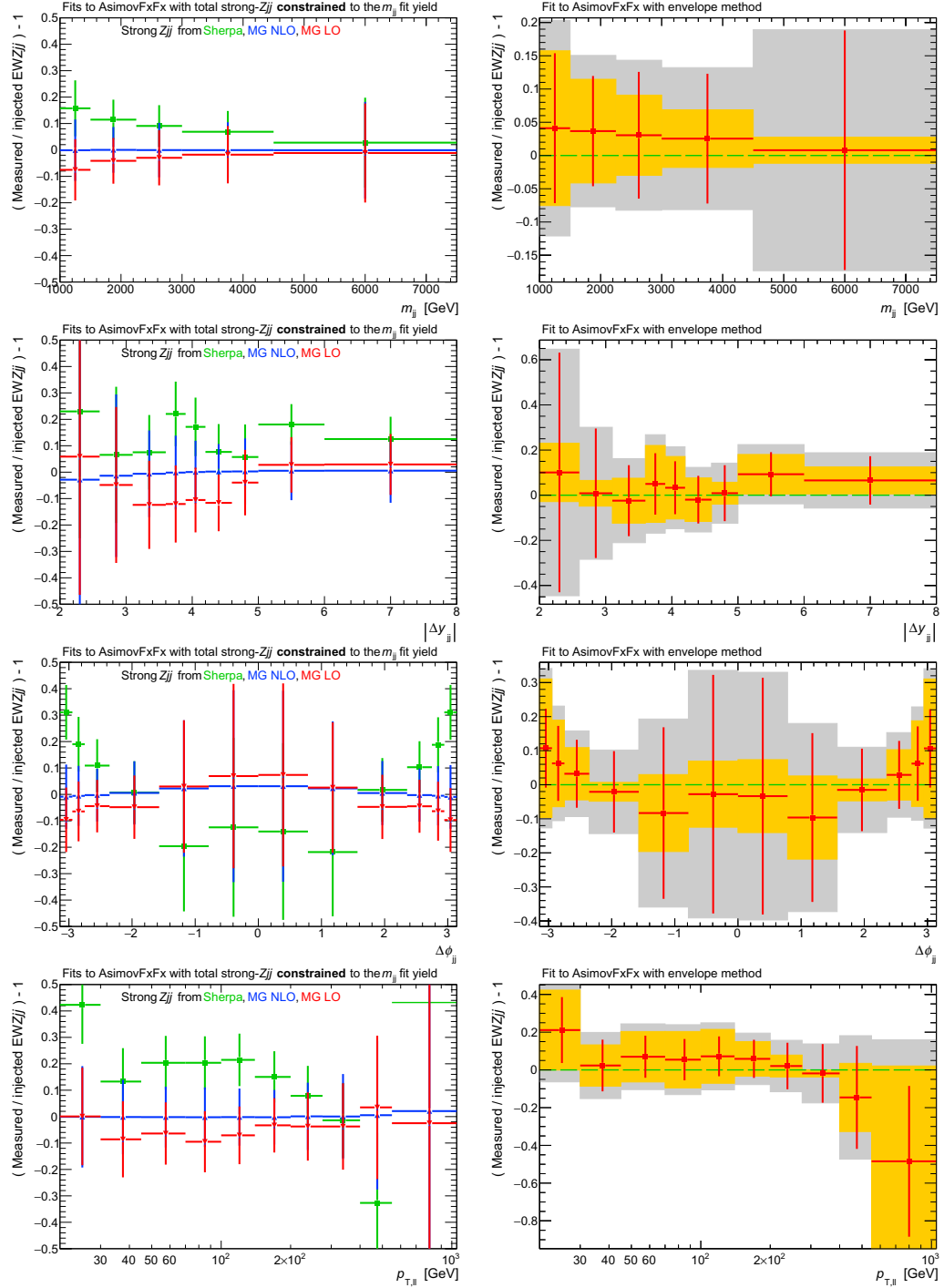


FIGURE 6.11: Measured EW- Zjj component for fits to AsimovFxFx with each strong- Zjj generator relative to the injected POWHEG+PY8 EW- Zjj prediction (left) and coverage of the strong- Zjj generator choice systematic (orange) from the envelope of the three EW extraction fits (right). The statistical uncertainty from the likelihood fit is shown as a gray band, this is the numerical average of the statistical uncertainty of each of the three fits.

6.4 Systematic uncertainties

The systematic uncertainties evaluated in the analysis fall into three categories. The first are the experimental systematics; these are mostly object associated to aspects of the determination of the leptons and jet objects including reconstruction, identification, isolation, jet calibration and jet flavour. Experimental systematics are typically evaluated as a set of uncorrelated sources of error that are referred to as *nuisance parameters* (NPs)*. The common analysis code used for object selection, provided by the various performance groups within the ATLAS collaboration, facilitates the evaluation of systematic NPs by modifying the properties of the relevant physics objects according to a 1σ variation of the systematic. The event selection is repeated after each unique systematic variation is applied, which will result in a modified set of events passing the selection and a migration of events between the signal and control regions. The impact of a particular NP is determined by comparing the systematically varied spectrum to the nominal spectrum. Besides the systematics associated to the jets and leptons there is also an experimental NP associated to the pileup reweighting procedure and uncertainty from the luminosity measurement.

The second category of systematic uncertainties are the theoretical uncertainties associated to the theoretical modelling implemented within the MC event generators. This includes the choice of the factorization (μ_f) and normalization (μ_r) scales as well as the uncertainty from the PDF set used in the simulation, as described in Section 2.3. The MADGRAPH generator was also used to evaluate the impact of the

*NP is used here as a shorthand to refer to an uncertainty source. The NPs correspond one to one with the uncertainty sources, the source has a set of process dependent uncertainty amplitudes δ_{ij} in each bin i that are controlled in the fit by an NP θ_j .

interference between EW- Zjj and strong- Zjj which is considered as an additional modelling NP of the EW- Zjj simulation.

The third category of systematic uncertainty is the modelling uncertainty associated to the analysis procedure itself. This category is dominated by the uncertainty associated to the choice of strong- Zjj generator. This is evaluated from the size of the envelope of the results obtained from the analysis performed with each of the three MC generators. There are also modelling systematic uncertainties related to the unfolding procedure detailed in Section 6.7.

Each experimental systematics and theoretical systematics NP is propagated through the EW- Zjj extraction procedure. This is accomplished by varying the MC *templates* used in the fit. The term MC template is used for a prediction from a given MC generator stored as a histogram that is used as input to the fit. The nominal (central value) result is obtained by extracting the EW signal with the nominal MC templates. Similarly, a varied result can be extracted by using systematically varied MC templates. The size of the systematic uncertainty assigned to a given NP is the difference between the nominal and varied results.

The NPs with small impact, defined as those which have an average impact of $< 1\%$ per bin, are introduced directly into the likelihood fit as in Equation (6.4). These NPs are said to be *profiled* within the likelihood fit, with θ_j representing the actual nuisance parameter within the fit. The uncertainties are treated as Gaussian such that a value $\theta_j = 1$ represents a 1σ pull on the i^{th} NP within the fit, the relevant MC template will be modified by a 1σ variation of the systematic as in Equation 6.5. The minimization will select the value of $\hat{\theta}_j$ that gives the best agreement with the data.

The modelling systematic associated to the choice of strong- Zjj generator is related to the method itself and are evaluated via a different procedure, *bootstrapping* [107]. The impact of the systematic is evaluated by creating a collection of pseudo-experiments where the dataset (either real data or Asimov) is uniquely varied. The collection of varied datasets approximate the distribution implied by the statistical uncertainty of the nominal dataset. The impact of the strong- Zjj generator systematic is evaluated by performing the EW- Zjj extraction procedure on each dataset in the statistically varied collection.

6.4.1 Experimental systematics

The majority of experimental systematic sources are object systematics related the physics objects used in the analysis; the jets, electrons and muons. The object systematics that were investigated will be enumerated in this section. These are the systematics recommended by the electron/photon, muon and jet/missing transverse energy combined performance groups. The other experimental systematic sources are the luminosity uncertainty of 1.7% [60] and the uncertainty from pileup reweighting. There is a nominal scale factor of 1/1.03 applied to reweight data to MC. This is to account for an effect in the MC that the μ distribution is too *hard*, meaning that a given MC events is best at simulating a data event with slightly higher μ [108]. There is a systematic uncertainty associated to this procedure, it is accounted for with an “up” and “down” variation of 1/0.94 and 1/1.12, respectively.

The following list of NP related to electrons was investigated:

- 1 energy scale NP
- 1 energy resolution NP

- 1 identification NP
- 1 isolation NP
- 1 reconstruction NP
- 2 trigger NPs

And, the following list of NPs related to muons was investigated:

- 1 momentum scale NP
- 2 momentum resolution NPs (1 in the ID, 1 in the MS)
- 2 identification NPs
- 2 isolation NPs
- 4 reconstruction NPs
- 2 trigger NPs
- 2 track to vertex association NPs
- 2 sagitta bias[†] NPs

A majority of these sources are found to have negligible (~ 0) impact, this is because the leptons selected in this analysis are required to have relatively high transverse momentum $p_T > 25 \text{ GeV}$ and η that falls within the inner detector. In general, the ATLAS detector measures leptons satisfying these requirements well and the systematic variations have little impact on the set of selected events. Three systematic NPs were found to have a non-negligible impact and thus were propagated through the analysis:

- 1 electron identification NP (EL_EFF_ID_TOTAL_1NPCOR_PLUS_UNCOR)

[†]*Sagitta* is a quantity that characterizes the curvature (momentum) of a particle track in a magnetic field. A sagitta bias is caused by a rotational misalignment of the detector layers, this bias will have an opposite effect on negatively and positively charged particles

- 1 muon trigger NP (MUON_EFF_TrigSystUncertainty)
- 1 muon sagitta bias NP (MUON_SAGITTA_RHO)

Systematic uncertainties associated to jet reconstruction are split into two broad categories, jet energy scale (JES) and jet energy resolution (JER). The complete list of JES/JER systematics is roughly 130 total nuisances parameters. Investigating each separately is only necessary for the most sensitive analyses, thus the jet performance group provides configurations that reduce the number of NPs. This analysis uses the so-called “category reduction” for JES systematics which reduces the number of NPs from roughly 100 to 29, leaving enough information that a combination with CMS is possible. For JER systematics, the “full JER” reduction scheme is used to reduce the number of NPs from 34 to 13. Additionally there is a single NP associated to the difference in efficiency between data and MC of the default JVT requirement that is applied as part of the jet selection.

The JES uncertainties are broken down by category as follows:

- 5 η intercalibration NPs (1 modelling, 3 non-closure, 1 statistical)
- 2 flavour NPs (flavour response and flavour composition)
- 1 b jet response NP
- 4 pileup NPs
- 1 punch through NP
- 1 extremely high p_T single particle NP
- 2 detector NPs
- 4 modelling NPs
- 6 statistical NPs
- 3 mixed NPs

JES uncertainties come from a variety of primary sources most of which are the jet calibration procedures such as η intercalibration, Z+jet balance and multi-jet balance. Other primary sources are MC non-closure, jet flavour and pileup. Each of the 29 NPs are evaluated and the ones with significant impact are propagated to the the electroweak extraction procedure, for JES the non-negligible NPs are:

- EtaIntercalibration_Modelling
- EffectiveNP_Modelling1
- Flavor_Composition
- Flavor_Response
- Pileup_OffsetMu
- Pileup_OffsetNPV
- Pileup_RhoTopology

The JER systematics with “full JER” reduction gives 13 NPs, 1 to account for data/MC discrepancies and 12 generic NPs reduced from primary JER sources. The 12 generic NPs produce variations of similar sizes (0-2%), thus each is propagated to the electroweak extraction. Note that since the insitu jet calibration is derived to large extent using γ +jet and $Z \rightarrow ee$ events, the systematic uncertainties associated with the scale are correlated with the electron scale uncertainties. However, since all insitu JES and the electron or muon scale uncertainties are found negligible, this is of no concern to the analysis.

Figures 6.12 and 6.13 show impact of four JES NPs and of the pileup reweighting (PRW) variation in each bin of the signal and control regions for each of the four measurement variables. These systematic sources have the largest impact on the

measurement yield on average. The PRW variation in particular impacts the number of gap jets and tends to cause a migration of events from the $N_{\text{jets}}^{\text{gap}} = 0$ regions to the $N_{\text{jets}}^{\text{gap}} \geq 1$ regions. The JES variations will cause events to migrate into or out of the VBF topology regions due to variations in jet p_T , this has the largest impact in CRa and CRb regions ($N_{\text{jets}}^{\text{gap}} \geq 1$) for EW- Zjj . This is because in EW- Zjj the gap jet multiplicity is smaller compared to strong- Zjj , when the JES is varied it is therefore more likely to migrate events into the phase space. In strong- Zjj , where gap jets are expected, it is therefore less likely for a variation to migrate events into the gap jet regions. This relatively large uncertainty in the control regions for EW- Zjj is not cause for concern because the EW- Zjj component in these regions is minimal, therefore even a relatively large fractional variation of $\sim 10\%$ will result in a small absolute variation compared to the size of the dominant strong- Zjj component in the control regions.

6.4.2 Theoretical systematics

The theoretical systematics discussed in this section were evaluated with a so-called Rivet [109] routine designed to replicate the analysis selection at particle-level. The particle-level selection is detailed in Section 5.3.

The variation of the QCD renormalization (μ_r) and factorization (μ_f) scale on the strong- Zjj process was studied using the SHERPA event generator. Each scale can be varied independently by a factor of two around the nominal scale choice, up and down variations of each scale are performed for each possible combination (up–down, down–up, up–up, etc.). The impact of 6 such combinations is shown in Figure 6.14, increasing the scales results in a lower cross section and hence less

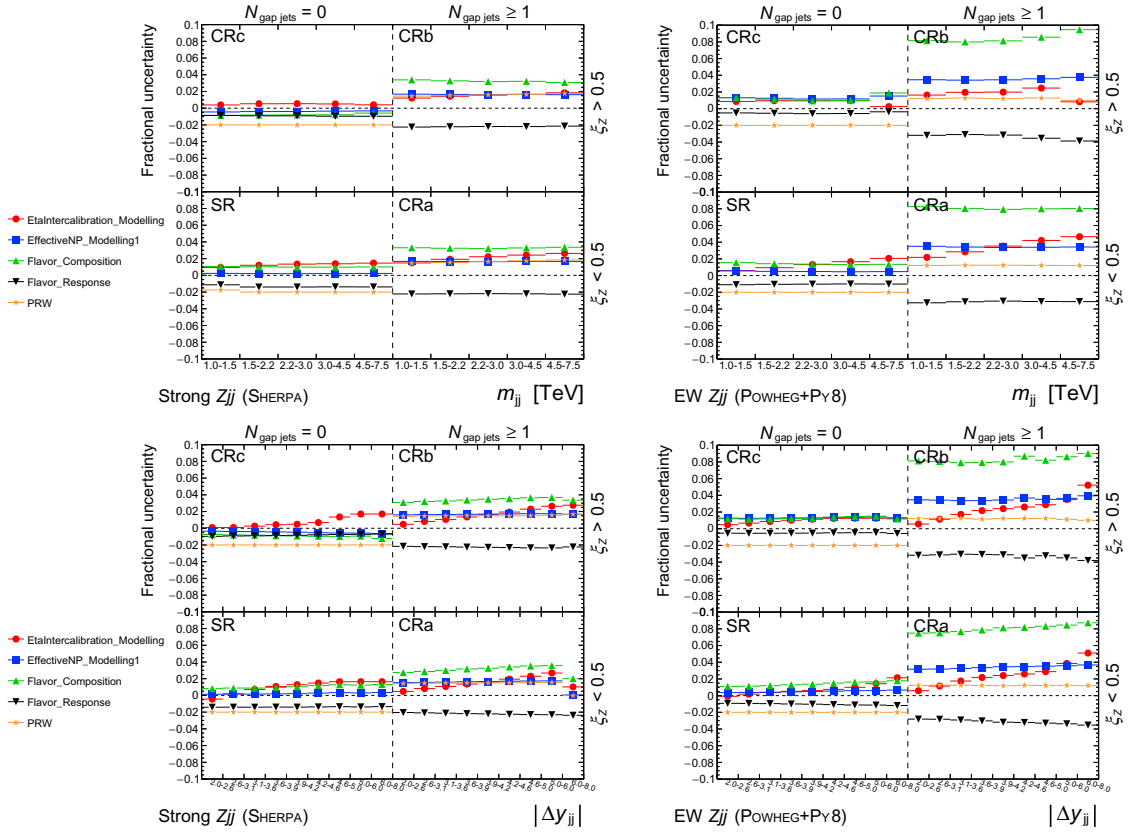


FIGURE 6.12: Impact of significant JES nuisance parameters and of the PRW variation as a function of m_{jj} (top) and $|\Delta y_{jj}|$ (bottom) in each of the control regions and in the signal region. The systematic uncertainties have been evaluated with the SHERPA strong- Zjj sample (left) and with the POWHEG+PY8 EW- Zjj sample (right). The large impact in the $N_{\text{jets}}^{\text{gap}} \geq 1$ regions for EW- Zjj is due to the smaller gap jet multiplicity that makes migrations into these regions more likely.

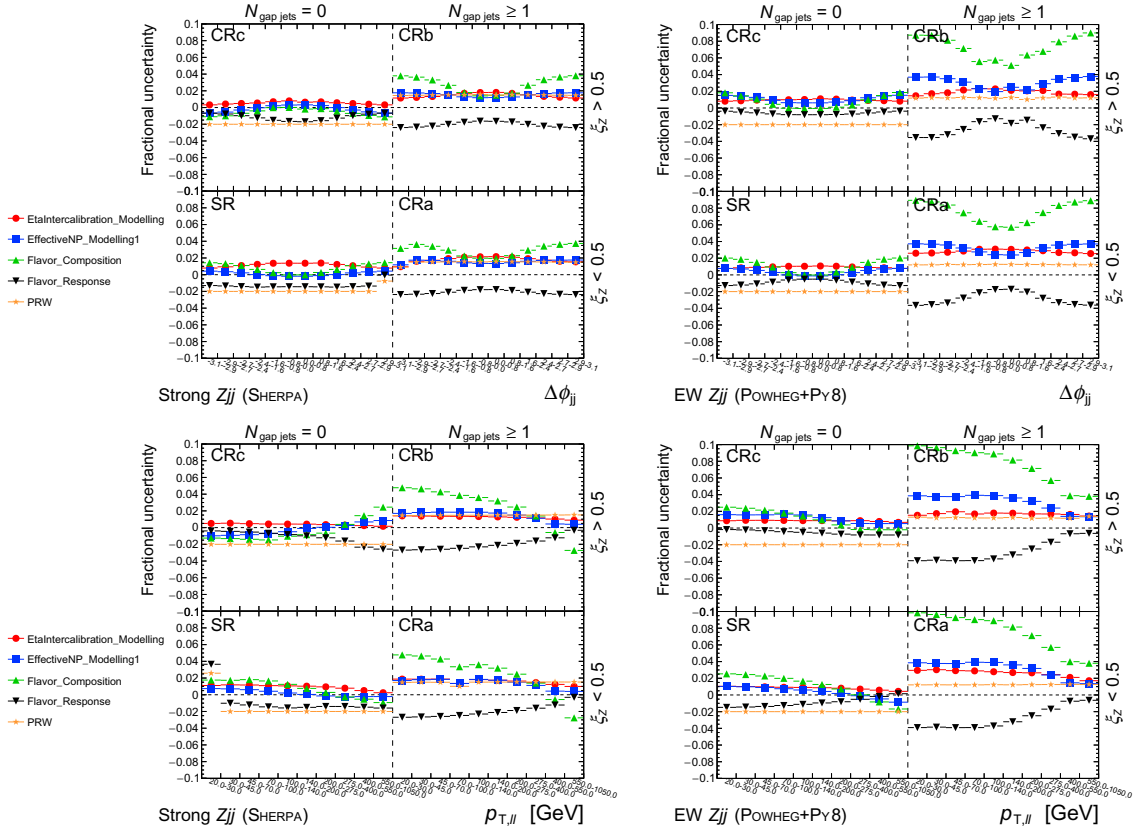


FIGURE 6.13: Impact of significant JES nuisance parameters and of the PRW variation as a function of $\Delta\phi_{jj}^{\text{signed}}$ (top) and $p_{T,\ell\ell}$ (bottom) in each of the control regions and in the signal region. The systematic uncertainties have been evaluated with the SHERPA strong- Zjj sample (left) and with the POWHEG+PY8 EW- Zjj sample (right). The large impact in the $N_{\text{gap jets}} \geq 1$ regions for EW- Zjj is due to the smaller gap jet multiplicity that makes migrations into these regions more likely.

events in the observed m_{jj} spectrum and vice versa. The renormalization scale variation is the dominant factor. The maximal up and down variations, those where both scales are varied in the same direction, are propagated through the analysis. Other combinations of μ_r and μ_f with produce variations that fall within the envelope of the two maximal variations.

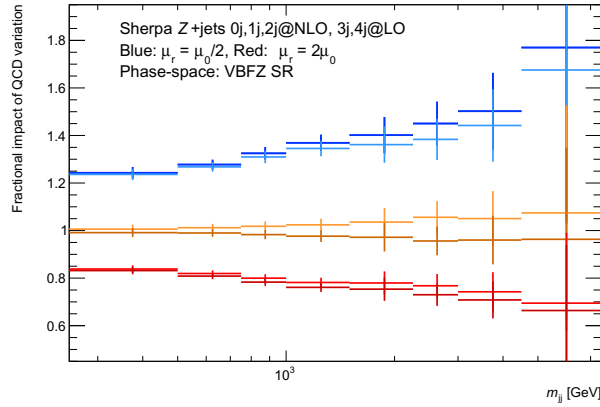


FIGURE 6.14: Impact of the QCD scale variations on the m_{jj} spectrum from SHERPA generated strong- Zjj events. The dark blue and dark red show the result of varying μ_r and μ_f in the same direction, down and up, respectively. The lighter blue and red show the result of only varying μ_r , and the yellow lines only varying μ_f .

The impact of the same scale variations, μ_r and μ_f , was also studied on the EW- Zjj process using the SHERPA event generator. This generator suffers from a known issue with the colour flow implementation (described in Section 4.1.2), however it is still the best available option for evaluating scale variations on the EW- Zjj process. As for strong- Zjj , only the two of the maximal variations of μ_r and μ_f are considered.

Scale variations were also studied for the HERWIG+VBFNLO generator for the EW- Zjj process. Both the hard scatter scale and the parton shower scale can be

varied up (2.0) or down (0.5) when creating the integration grids from which particle-level events are generated. By observing the effect these scale variations have on the EW- Zjj cross section, the size of the scale systematic was estimated. Averaging the impact of the up and down variations, the parton shower scale is 8.5% and the hard scale is 3.0%, together giving an estimate of 9.0% on the EW- Zjj scale.

The variability of the underlying PDF used in the strong- Zjj and EW- Zjj generation was also studied using the SHERPA event generator. The nominal PDF used by the generator is the central value of some underlying PDF distribution, the nominal PDF is provided by NNPDF3.0 [73] and interfaced to SHERPA and 100 variations to this nominal PDF are also provided. The procedure outlined by NNPDF is followed to evaluate the uncertainty. The RMS of the distribution of these 100 variations is assigned as the 1σ variation of the PDF uncertainty. Some bins suffer from large variations due to single events with large weights, this is not a physical effect, and was accounted for by applying a smooth variation when implementing this systematic variation. This PDF variation has a much smaller impact of typically 1-2% compared to the large QCD scale variation. A single up and down variation of the α_s value used in the PDF was also evaluated using the SHERPA event generator. The size of this α_s variation is 0-2% in EW- Zjj and 1-5% in strong- Zjj , thus it is also dominated by the QCD scale uncertainty.

Figures 6.15 and 6.16 show impact of the QCD scale and PDF variations in each bin of the four measurement variables. The size of the QCD scale variation is typically larger in strong- Zjj than EW- Zjj , particularly when focussing on the SR where the majority of EW- Zjj events fall.

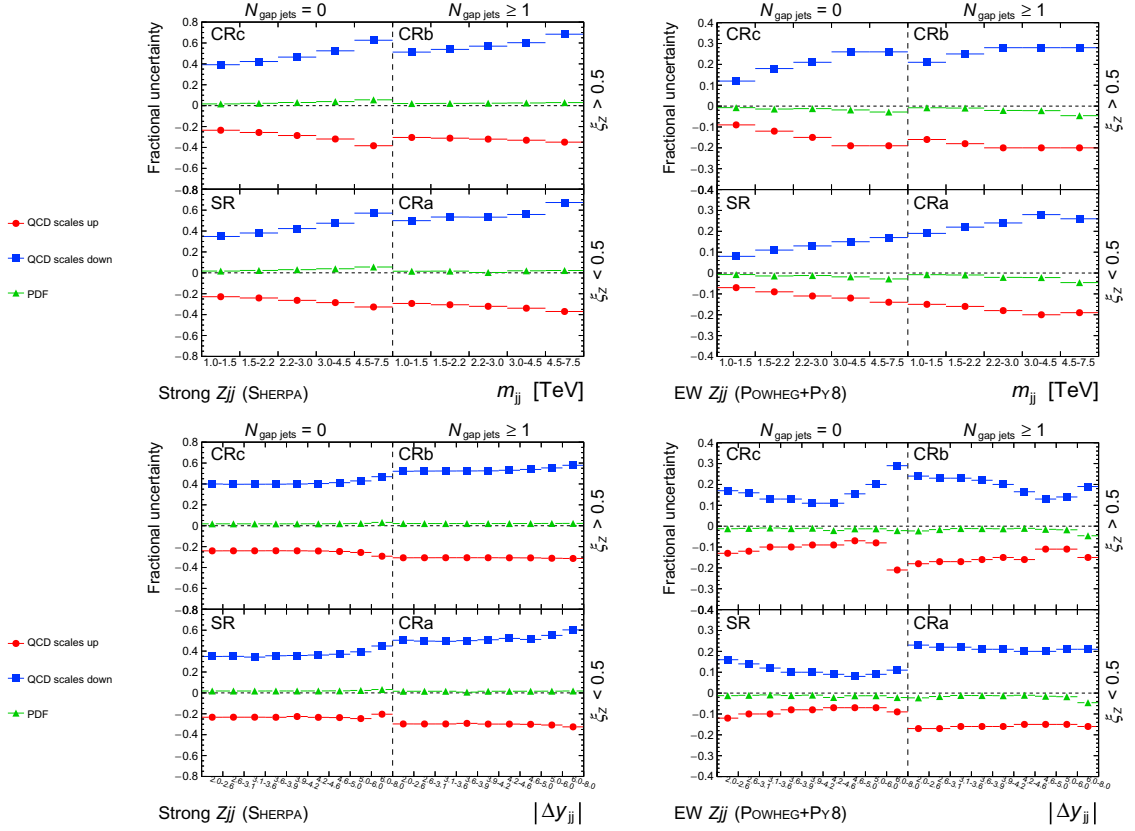


FIGURE 6.15: Impact of QCD scale and PDF variations as a function of m_{jj} (top) and $|\Delta y_{jj}|$ (bottom) in each of the control regions and in the signal region. The systematic uncertainties have been evaluated with the SHERPA strong- Zjj sample (left) and with the POWHEG+PY8 EW- Zjj sample (right). Note that the range of the y-axis is different for the left and right figures. The QCD scale variations impact EW- Zjj less than strong- Zjj as the former process has fewer QCD vertices.

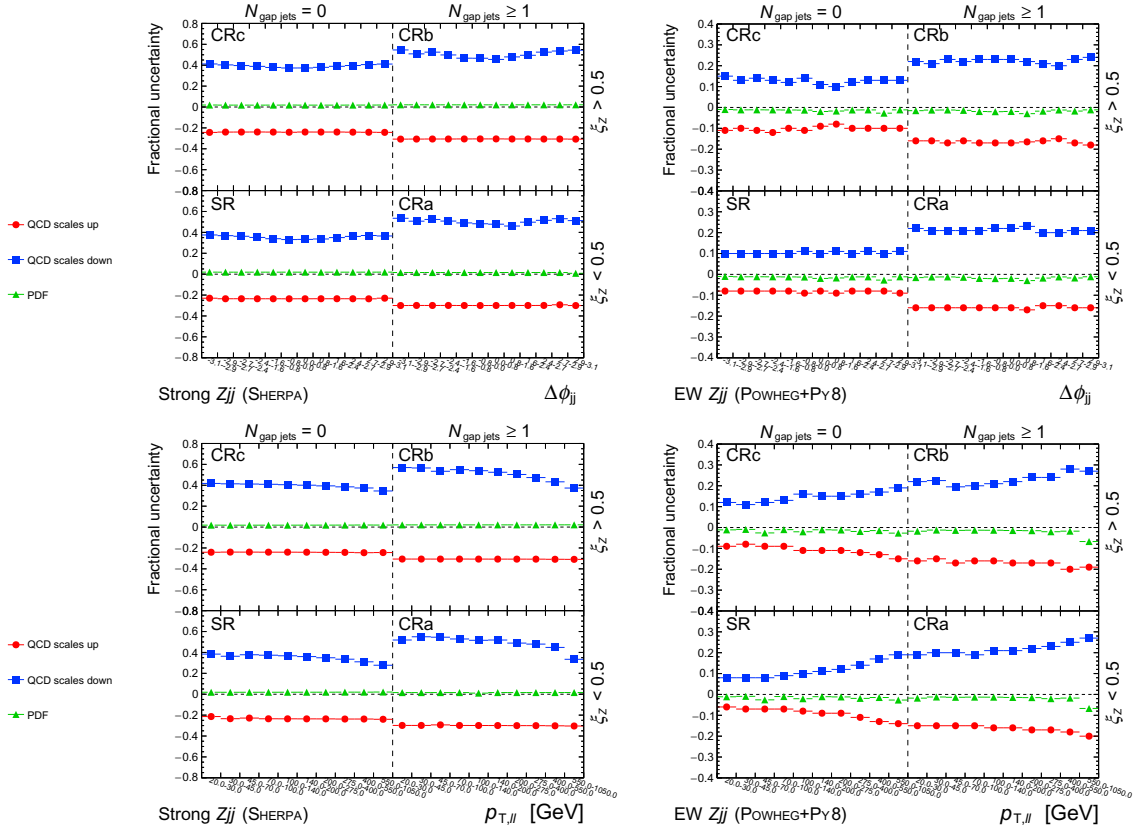


FIGURE 6.16: Impact of QCD scale and PDF variations as a function of $\Delta\phi_{jj}^{\text{signed}}$ (top) and $p_T^{\ell\ell}$ (bottom) in each of the control regions and in the signal region. The systematic uncertainties have been evaluated with the SHERPA strong- Zjj sample (left) and with the POWHEG+PY8 EW- Zjj sample (right). Note that the range of the y-axis is different for the left and right figures. The QCD scale variations impact EW- Zjj less than strong- Zjj as the former process has fewer QCD vertices.

The interference between EW- Zjj and strong- Zjj discussed in Section 5.2 is evaluated with the same `Rivet` routine used to evaluate the other theory systematics. This interference is treated as an additional systematic on the EW- Zjj signal and propagated through the analysis. The impact of the interference is between 1–4% in a typical bin of the measurement.

6.4.3 Bootstrapping the EW- Zjj extraction

As described in Section 6.1, the EW- Zjj component is extracted three times, once using the SHERPA strong- Zjj prediction, once using the MG5+PY8 strong- Zjj and once using the MG5_NLO+PY8 reweighted strong- Zjj prediction. The final electroweak signal yield in each bin of the differential distribution is taken to be the midpoint of the *envelope* of yields obtained using the three different strong- Zjj event generators. Here the envelope means the difference between the largest and smallest from the set of the three strong- Zjj generator results, this is illustrated in Figure 6.20. The size of the envelope defines the systematic uncertainty associated with the choice of strong generator. This source is typically the largest systematic uncertainty in any particular measurement bin. To better understand the behaviour of this uncertainty it is estimated via bootstrapping with pseudo-experiments. With this technique one can simulate the underlying statistical distribution that the nominal (central value) result is sampled from. This method is also useful for evaluating the statistical uncertainty of the EW- Zjj extraction with the spread of the pseudo-experiment results.

For each data (real or simulated) event that is part of the signal or control regions, a random number is generated according to a Poisson distribution with mean equal

to the event weight from MC or equal to 1 for data. Then the four measurement distributions are populated according to this randomly generated Poisson weight. For this study 500 such datasets were generated from both real data and for two Asimov datasets, AsimovFxFx and AsimovSrew described in Section 6.3. The 500 datasets are fit individually, this is considered as a pseudo-experiment. Quantities such as the generator choice systematic are computed from the results of each pseudo-experiment.

The sign of the generator uncertainty is assigned by observing the fit to SHERPA strong- Zjj relative to the centre of the three fit envelope. The SHERPA typically forms one edge of the envelope and gives a different prediction relative to the two MG5+PY8 based fits.

The evaluated systematic and strong- Zjj generator choice systematic from the fits of the real data pseudo-experiments are shown in Figure 6.17. Note that the object and theoretical systematic uncertainty values presented thus far were the pre-fit values, *i.e.* the variation observed in the MC templates input to the likelihood fit. The systematics evaluated with bootstrapping contribute to the final uncertainty on the extracted EW- Zjj result that will be presented in Section 6.5.

Besides evaluating these uncertainties, the bootstrap method was used to further validate the method with Asimov data. The EW- Zjj component of the simulated data is known, it is equal to the nominal POWHEG+PY8 prediction. By performing the pseudo-experiments and extracting an EW- Zjj component for each dataset, one can study the performance of the method by comparing the measured EW component to the known (*injected*) prediction. Figure 6.18 shows the mean and RMS of the distribution of measured / injected EW- Zjj from the 500 pseudo-experiments.

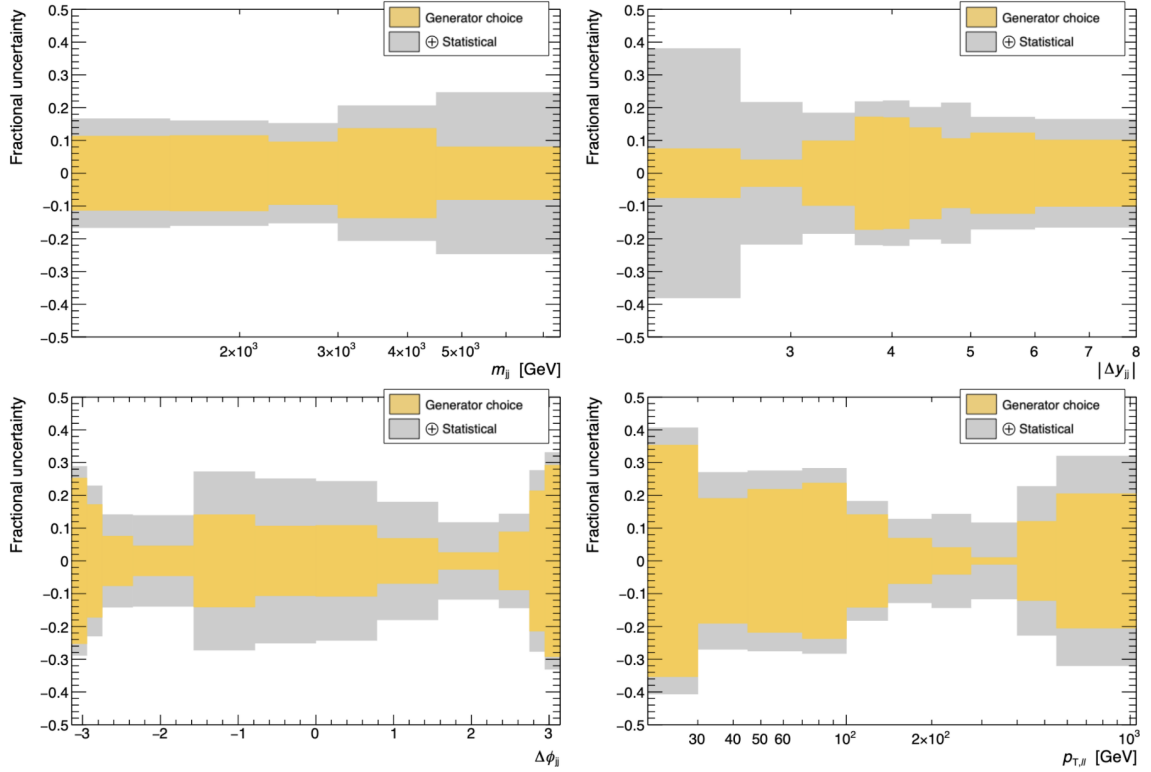


FIGURE 6.17: Statistical and strong- Zjj generator choice uncertainties evaluated from the bootstrapped real data datasets. The grey band shows the two uncertainties added in quadrature.

The generator choice systematic is a crucial uncertainty of the measurement. This bootstrap study provides a more detailed understanding of the behaviour of this important uncertainty on average. The figure shows that the spread of the bootstrapped generator choice systematic is consistent with the injected EW- Zjj value in the majority of bins for both AsimovFxFx and AsimovSrew.

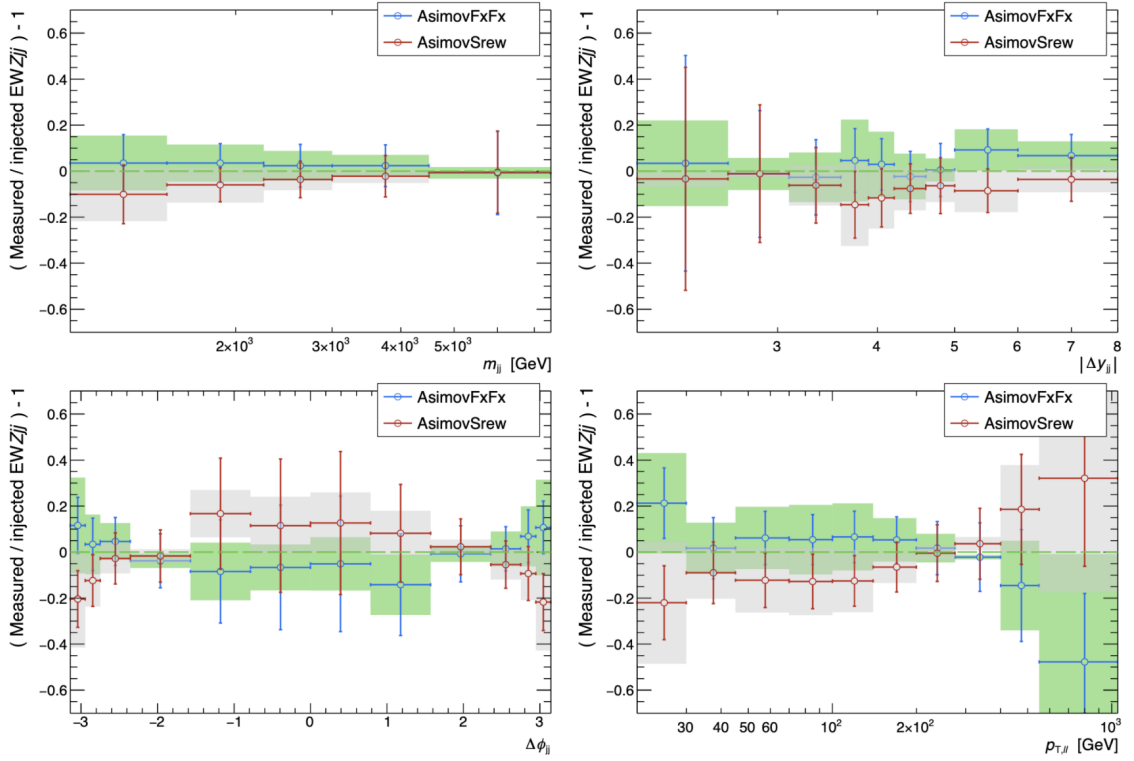


FIGURE 6.18: Comparison of measured to injected EW-Z jj signal from the fits of the bootstrapped Asimov datasets. The error bars show the spread of the measured EW-Z jj component from the 500 pseudo-experiments and the error band is the spread of the measured generator systematics.

6.5 Detector-level EW- Zjj results

With all sources of systematic uncertainty evaluated, the likelihood fit outlined in Section 6.1 was performed to extract the EW- Zjj component from the real ATLAS data. The pre-fit and post-fit agreement between data and simulation is shown in Figure 6.19 as a function of m_{jj} in the signal and control regions. Two separate fits are shown, one using the SHERPA strong- Zjj prediction (top row) and one using the MG5_NLO+PY8 prediction (bottom row). These simulations initially have very different mismodelling as a function of m_{jj} , but produce very good agreement with the data following the fitting procedure. The values of the fitted parameter for each measurement variable and for the fits with each of the three strong- Zjj predictions are listed in Tables 6.2–6.5.

The centre of envelope formed by the results of the three fits defines the central value of the nominal result and the size of the envelope determines the generator choice systematic, this is depicted in Figure 6.20 for each bin of the full measurement. Each measurement variable is presented separately in Figures 6.21 and 6.22. These show the detector-level differential cross sections and the breakdown of the measurement uncertainty in terms of statistical uncertainty and categories of systematic sources.

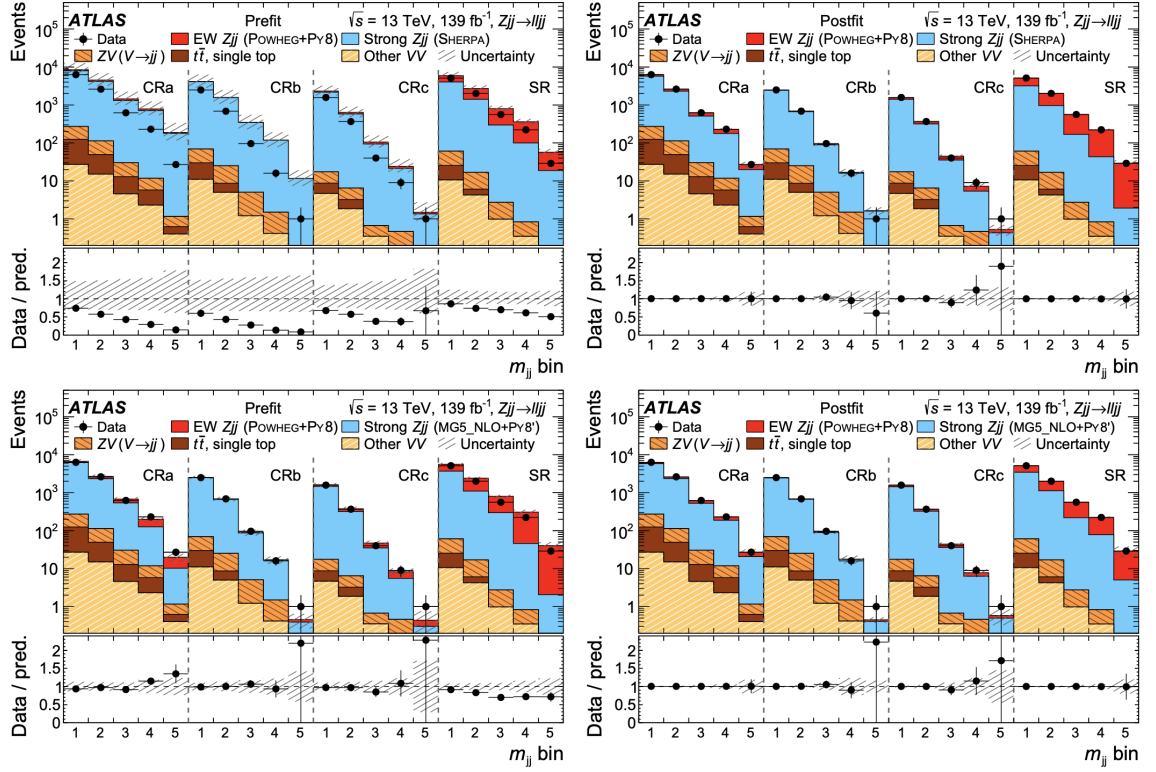


FIGURE 6.19: Comparison of the predicted and observed event yields before (left) and after (right) the likelihood fit using strong- Zjj estimates from SHERPA (top) and MG5_NLO+PY8' (bottom). Before the fit the EW- Zjj component is the POWHEG+PY8 prediction and after the fit the EW- Zjj component is the measured yield. The m_{jj} bin edges are defined by (1.0, 1.5, 2.25, 3.0, 4.5, 7.5) TeV.

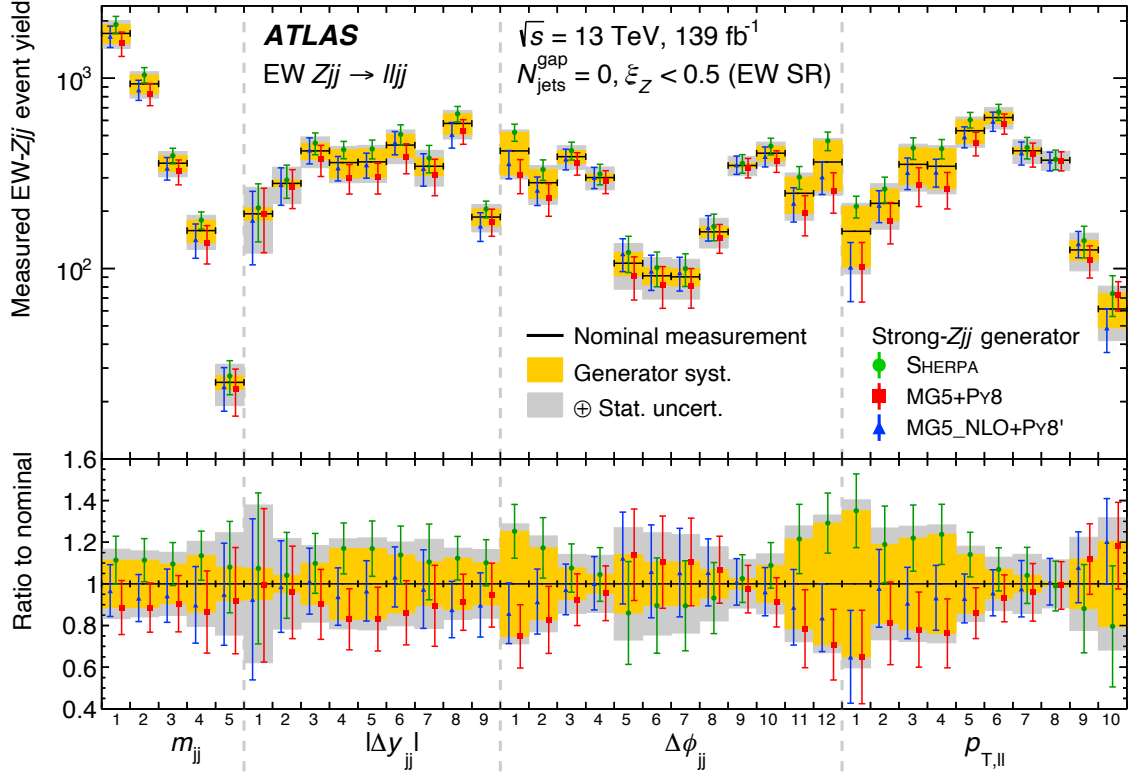


FIGURE 6.20: Measured EW- Zjj signal from the nominal method with the result from each strong- Zjj generator overlaid. The error bands correspond to the nominal statistical and generator systematic uncertainties, the error bars show the statistical uncertainty of the result from each generator. The generator systematic band is defined by the maximum and minimum values from the three individual results.

SHERPA strong- Zjj					
$b_{H,i}$	0.59	0.42	0.25	0.13	0.13
$b_{L,i}$	0.71	0.54	0.37	0.24	0.11
$f(x_i) \cdot b_{L,i}$	0.79	0.71	0.56	0.43	0.26
μ^{EW}	1.00	0.79	0.77	0.68	0.71
MG5+PY8 strong- Zjj					
$b_{H,i}$	0.84	0.74	0.60	0.52	0.52
$b_{L,i}$	0.75	0.74	0.71	0.95	1.16
$f(x_i) \cdot b_{L,i}$	0.53	0.54	0.53	0.74	0.96
μ^{EW}	0.80	0.63	0.64	0.52	0.61
MG5_NLO+PY8' strong- Zjj					
$b_{H,i}$	0.99	1.02	0.99	1.08	1.08
$b_{L,i}$	0.95	1.02	1.04	1.55	2.12
$f(x_i) \cdot b_{L,i}$	0.94	1.04	1.08	1.68	2.48
μ^{EW}	0.87	0.67	0.67	0.55	0.63
m_{jj} low bin edge [TeV]	1.0	1.5	2.2	3.0	4.5
m_{jj} high bin edge [TeV]	1.5	2.2	3.0	4.5	7.5

TABLE 6.2: Table of parameters values determined from the m_{jj} fit from each of the three strong- Zjj generators. A perfect strong- Zjj generator would result in $b_i = 1$ and $f(x_i) = 1$, MG5_NLO+PY8 performs the best in this respect.

SHERPA strong- Zjj									
$b_{H,i}$	0.60	0.56	0.54	0.48	0.50	0.45	0.46	0.38	0.38
$b_{L,i}$	0.70	0.63	0.63	0.59	0.65	0.62	0.63	0.49	0.39
$f(x_i) \cdot b_{L,i}$	0.79	0.74	0.74	0.71	0.79	0.77	0.80	0.64	0.54
μ^{EW}	1.42	1.21	1.06	1.13	0.95	0.78	0.63	0.75	0.78
MG5+Py8 strong- Zjj									
$b_{H,i}$	0.68	0.80	0.80	0.81	0.87	0.84	1.08	0.96	0.96
$b_{L,i}$	0.58	0.57	0.65	0.70	0.79	0.82	0.90	0.89	1.29
$f(x_i) \cdot b_{L,i}$	0.45	0.43	0.47	0.48	0.53	0.53	0.57	0.52	0.64
μ^{EW}	1.33	1.12	0.88	0.80	0.67	0.59	0.51	0.61	0.67
MG5_NLO+Py8' strong- Zjj									
$b_{H,i}$	0.95	1.03	0.96	0.95	1.00	0.96	1.21	1.04	1.04
$b_{L,i}$	0.92	0.83	0.87	0.91	1.01	1.00	1.10	1.08	1.46
$f(x_i) \cdot b_{L,i}$	0.96	0.84	0.86	0.88	0.96	0.94	1.01	0.95	1.18
μ^{EW}	1.22	1.15	0.98	0.92	0.78	0.71	0.56	0.58	0.64
$ \Delta y_{jj} $ low bin edge	2.0	2.6	3.1	3.6	3.9	4.2	4.6	5.0	6.0
$ \Delta y_{jj} $ high bin edge	2.6	3.1	3.6	3.9	4.2	4.6	5.0	6.0	8.0

TABLE 6.3: Table of parameters values determined from the $|\Delta y_{jj}|$ fit from each of the three strong- Zjj generators.

SHERPA strong- Zjj														
$b_{H,i}$	0.45	0.53	0.56	0.59	0.57	0.49	0.49	0.57	0.59	0.56	0.53	0.45		
$b_{L,i}$	0.53	0.59	0.61	0.63	0.67	0.63	0.63	0.67	0.63	0.61	0.59	0.53		
$f(x_i) \cdot b_{L,i}$	0.68	0.74	0.75	0.72	0.71	0.61	0.61	0.71	0.72	0.75	0.74	0.68		
μ^{EW}	1.00	0.79	0.86	0.83	0.67	1.08	1.07	1.06	0.94	0.91	0.72	0.90		
MG5+Py8 strong- Zjj														
$b_{H,i}$	0.80	0.89	0.81	0.75	0.72	0.83	0.83	0.72	0.75	0.81	0.89	0.80		
$b_{L,i}$	0.72	0.73	0.73	0.78	0.83	0.84	0.84	0.83	0.78	0.73	0.73	0.72		
$f(x_i) \cdot b_{L,i}$	0.59	0.57	0.53	0.48	0.38	0.26	0.26	0.38	0.48	0.53	0.57	0.59		
μ^{EW}	0.60	0.56	0.74	0.76	0.89	1.33	1.32	1.22	0.90	0.76	0.47	0.50		
MG5_NLO+Py8' strong- Zjj														
$b_{H,i}$	1.03	1.11	0.99	0.89	0.83	0.91	0.91	0.83	0.89	0.99	1.11	1.03		
$b_{L,i}$	0.99	0.98	0.96	0.97	1.02	1.01	1.01	1.02	0.97	0.96	0.98	0.99		
$f(x_i) \cdot b_{L,i}$	1.11	1.06	0.96	0.83	0.66	0.46	0.46	0.66	0.83	0.96	1.06	1.11		
μ^{EW}	0.69	0.62	0.78	0.81	0.88	1.31	1.29	1.22	0.94	0.80	0.52	0.59		
$\Delta\phi_{jj}^{\text{signed}}$ low bin edge [$\pi/16$]	-16	-15	-14	-12	-8.0	-4.0	0.00	4.0	8.0	12	14	15		
$\Delta\phi_{jj}^{\text{signed}}$ high bin edge [$\pi/16$]	-15	-14	-12	-8.0	-4.0	0.00	4.0	8.0	12	14	15	16		

TABLE 6.4: Table of parameters values determined from the $\Delta\phi_{jj}^{\text{signed}}$ fit from each of the three strong- Zjj generators.

SHERPA strong- Zjj										
$b_{H,i}$	0.45	0.48	0.47	0.53	0.47	0.52	0.57	0.61	0.61	0.78
$b_{L,i}$	0.59	0.62	0.58	0.58	0.55	0.61	0.64	0.61	0.65	0.64
$f(x_i) \cdot b_{L,i}$	0.75	0.78	0.73	0.72	0.68	0.75	0.77	0.72	0.74	0.67
μ^{EW}	1.07	0.75	0.76	0.72	0.96	0.97	0.87	1.05	0.93	1.14
MG5+PY8 strong- Zjj										
$b_{H,i}$	1.01	1.04	0.90	0.91	0.79	0.76	0.69	0.66	0.69	0.73
$b_{L,i}$	0.93	0.84	0.82	0.77	0.72	0.76	0.74	0.64	0.57	0.42
$f(x_i) \cdot b_{L,i}$	0.75	0.67	0.64	0.59	0.54	0.54	0.50	0.40	0.31	0.15
μ^{EW}	0.52	0.51	0.49	0.45	0.73	0.84	0.81	1.06	1.17	1.68
MG5_NLO+PY8' strong- Zjj										
$b_{H,i}$	1.28	1.31	1.14	1.14	0.98	0.94	0.83	0.79	0.84	0.87
$b_{L,i}$	1.22	1.08	1.07	0.99	0.94	0.99	0.98	0.84	0.77	0.57
$f(x_i) \cdot b_{L,i}$	1.38	1.21	1.18	1.06	0.98	1.00	0.92	0.72	0.56	0.25
μ^{EW}	0.52	0.62	0.57	0.54	0.79	0.87	0.83	1.07	1.15	1.70
$p_T^{\ell\ell}$ low bin edge [GeV]	20	30	45	70	100	140	200	275	400	550
$p_T^{\ell\ell}$ high bin edge [GeV]	30	45	70	100	140	200	275	400	550	1050

TABLE 6.5: Table of parameters values determined from the $p_T^{\ell\ell}$ fit from each of the three strong- Zjj generators.

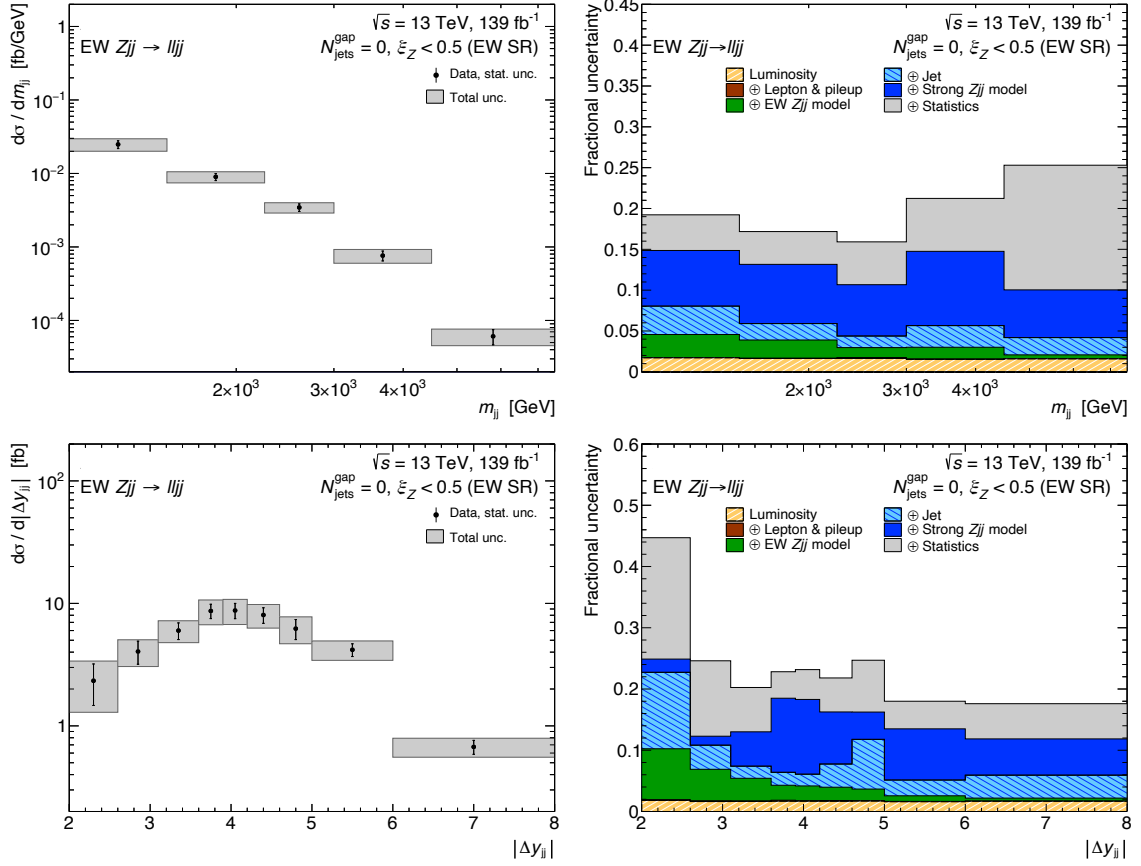


FIGURE 6.21: Differential cross-sections at detector-level for EW-Z jj production as a function m_{jj} (top) and $|\Delta y_{jj}|$ (bottom). The fractional uncertainty on the measurement is also shown, uncertainty sources are grouped in categories that are added in quadrature (denoted $\oplus$) to give the total uncertainty.

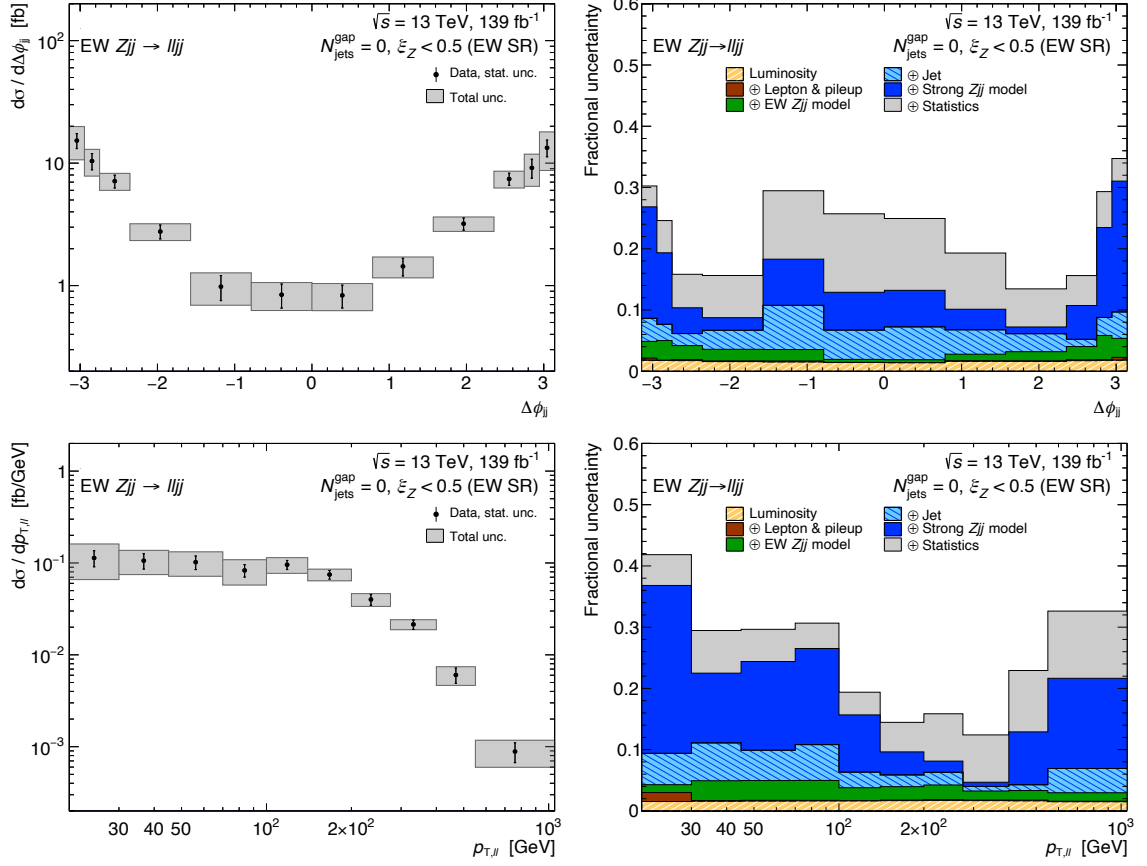


FIGURE 6.22: Differential cross-sections at detector-level for EW-Zjj production as a function $\Delta\phi_{jj}^{\text{signed}}$ (top) and $p_T^{\ell\ell}$ (bottom). The fractional uncertainty on the measurement is also shown, uncertainty sources are grouped in categories that are added in quadrature (denoted $\oplus$) to give the total uncertainty.

6.6 Statistical method variations and sequential alternative

The nominal results presented in the previous section were obtained with the nominal statistical method to extract the EW- Zjj component. The nominal method has three key features that can be varied. These are: it uses CRA to directly constrain the SR with the bin dependent factors; it uses a two parameter linear function to apply the residual correction; and, it also applies the additional integral constraint on the strong- Zjj component in the SR such that the total yield matches between the four measurement variables. This section will show how the results are effected by taking an alternative approach for each of these features. These are: using CRC instead to directly constrain the SR (this alternative was presented in Section 6.2.1); introducing an additional degree of freedom to the residual correction function; and, removing the additional constraint on the strong- Zjj yield.

In addition to these variations to the nominal maximum likelihood method, a separate alternative approach was used to extract the EW- Zjj component. This alternative will be referred to as the *sequential* method, so-called because the constraints applied to correct the strong- Zjj component are derived and applied sequentially. This method does not make use of a likelihood function, instead constraint functions are fit to ratios of the data and MC predictions between the control and signal regions. In the following sections each of these four alternatives will be described and the results obtained with each will be presented.

6.6.1 Orthogonal control regions

This variation consists of exchanging the roles of CRa and CRc in the model that constrains the SR (Equation (6.3)). The bin dependent factors are now $b_{\text{NGJ},i}$ (no gap jets) applied in CRc and in the SR and $b_{\text{GJ},i}$ (gap jets) applied in CRa and CRb. The residual correction function is applied in CRa and in the SR and the model is otherwise unchanged. As depicted in Figure 6.23, this variation can be thought of as a rotation of the constraints by $\pi/2$ in the $N_{\text{jets}}^{\text{gap}} - \xi_Z$ space in which the signal and control regions are defined. Recall that in Section 6.2.1 it was shown that this variation is expected to yield similar results to the nominal method. The choice to use CRa for the bin dependent factors shared with the SR as the nominal method is made because CRa has more statistics than CRc, particularly in the last bin of m_{jj} and $|\Delta y_{jj}|$.

The results from the variation are compared to the nominal detector-level result in Figure 6.24. The statistical and generator choice systematic uncertainties for the varied case were determined with bootstraps as is done in the nominal method. It is observed that the nominal result falls within the uncertainty of the varied method in all bins. A formal compatibility test was performed, the results of these tests for all the alternative methods is summarized in Section 6.6.5.

6.6.2 Functional form variation

A two parameter linear function $f(x) = a|x| + b$ is used in the nominal method to constrain the strong- Zjj component in CRc and in the SR. This variation introduces a third parameter such that $f(x) = ax^2 + b|x| + c$. This variation corresponds to a particular choice of a polynomial parameterization. The results from the variation

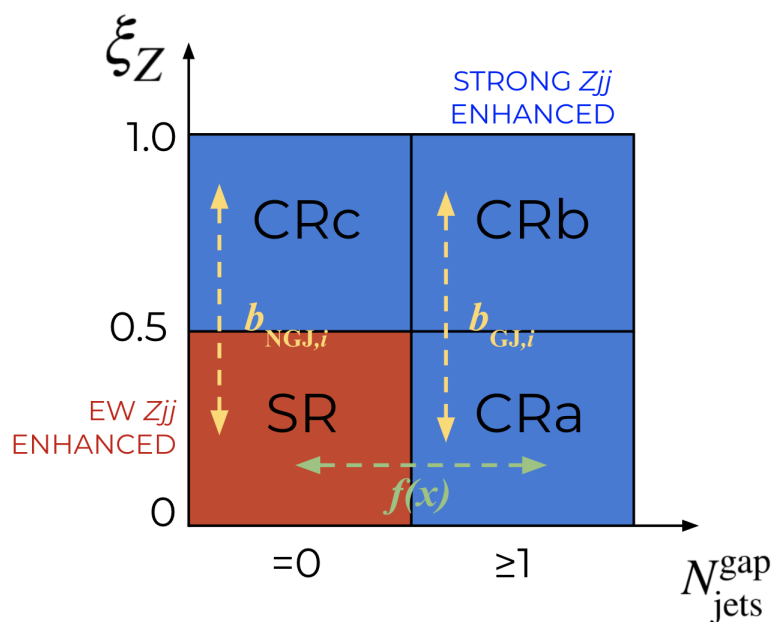


FIGURE 6.23: Diagram depicting the alternative CRc method. The $b_{\text{NGJ},i} / b_{\text{GJ},i}$ factors and the $f(x)$ function are applied to the signal and control regions to constrain the strong- Zjj component. These constraints have been in a sense rotated by $\pi/2$ with respect to those in the nominal method.

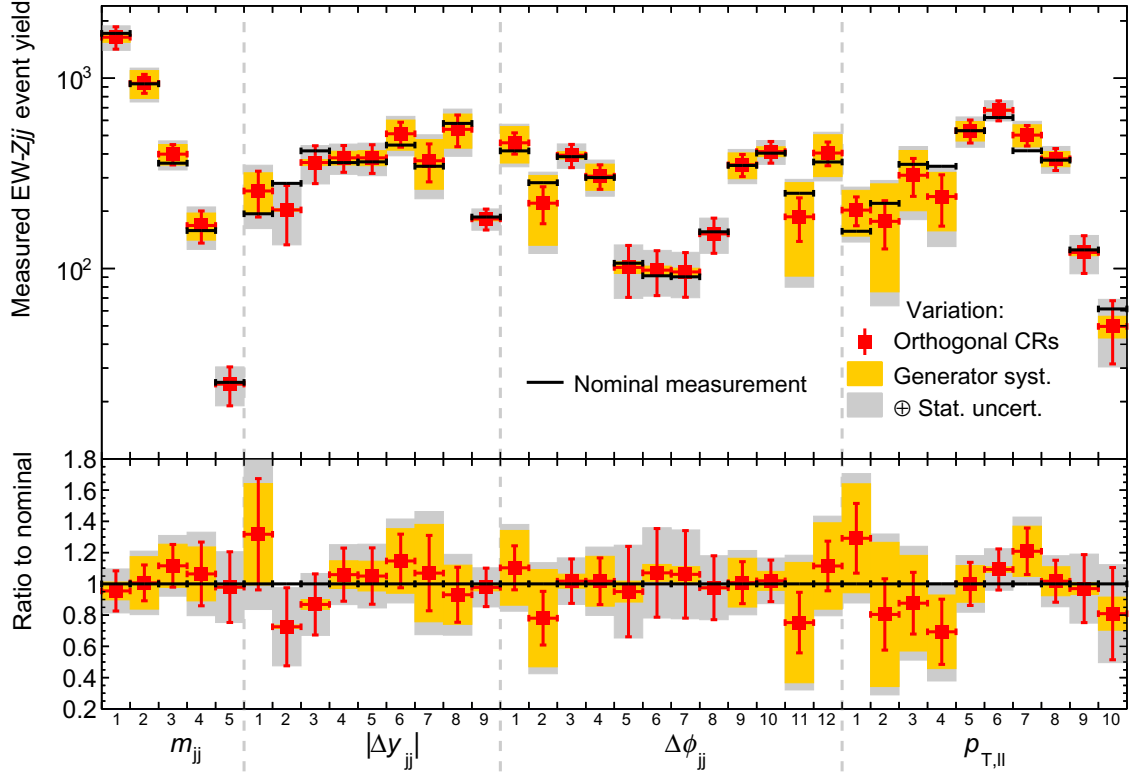


FIGURE 6.24: Measured EW-Zjj signal from the varied method using orthogonal control regions. The statistical uncertainty is shown as an error bar, the strong generator choice systematic is shown in orange and the combined statistical plus generator uncertainty is shown in grey. All uncertainties are derived from bootstraps and are from the varied method, the central value of the nominal result is shown for comparison.

are compared to the nominal detector-level result in Figure 6.25. Adding the extra degree of freedom to the fit function produces a result that is in good agreement with the nominal measurement in all bins. The two parameter function was therefore found to be sufficient. The flexibility of the higher order function is more sensitive to statistical variations of the bootstraps, particularly in the bins at the edges of the distributions leading to inflated statistical uncertainty. Alternative two parameter functions such as $f(x) = ae^{b|x|}$ (exponential) and $f(x) = a|x|^b$ (power law) were also studied and found to produce similar results compatible with the nominal linear function.

Overall, the likelihood fit method is found to not be particularly sensitive to the choice of functional form, thus a linear parameterization was chosen for its relative simplicity and for the stability of the fitted results in the bootstrapped pseudo-experiments.

6.6.3 Free strong- Zjj yield

The constraint on the integral of the corrected strong- Zjj prediction in the signal region is implemented to ensure consistency between the four measurement variables. As described by Equation (6.6), the strong prediction in the non- m_{jj} variables is constrained to equal the yield observed in the m_{jj} measurement. The m_{jj} variable is used for this purpose because the EW and strong Zjj processes have very different shapes as a function of m_{jj} and the measurement has the fewest bins, this result in the most precise measurement of the integrated strong- Zjj event yield.

This variation removes the constraint, as such the non- m_{jj} measurements may converge to different overall strong- Zjj event yields. The results from the variation

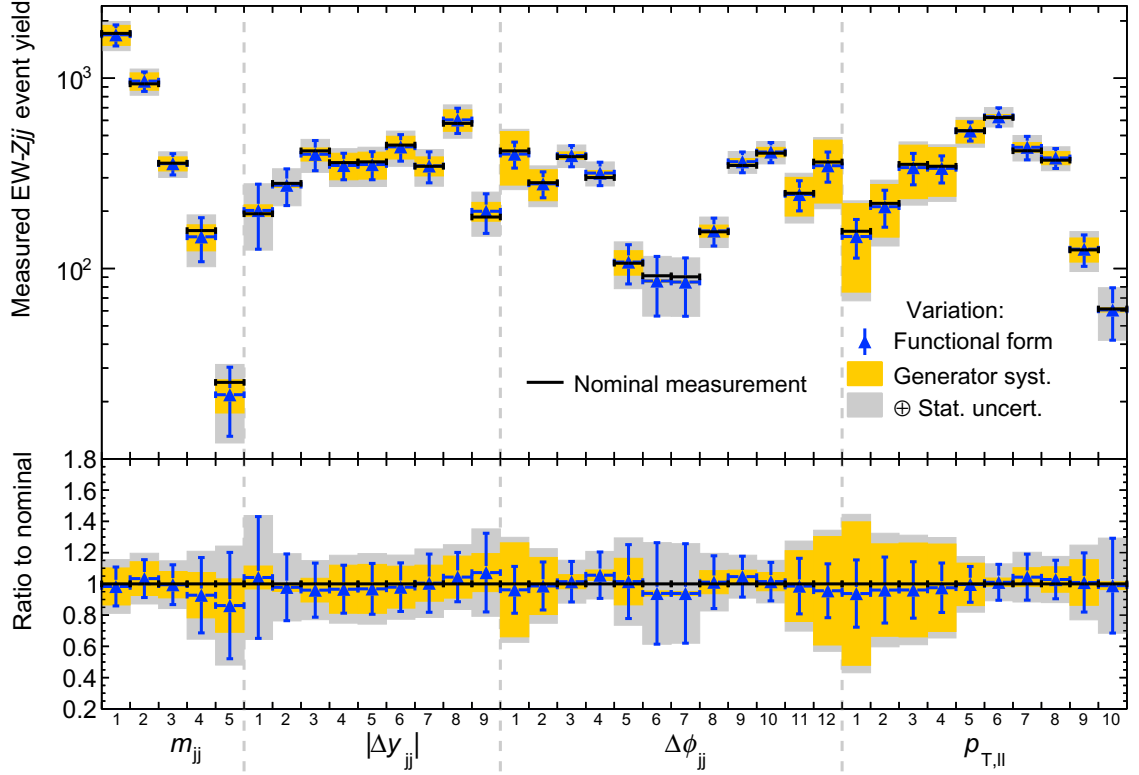


FIGURE 6.25: Measured EW- Zjj signal from the varied method using a degree two polynomial fit function. The statistical uncertainty is shown as an error bar, the strong generator choice systematic is shown in orange and the combined statistical plus generator uncertainty is shown in grey. All uncertainties are derived from bootstraps and are from the varied method, the central value of the nominal result is shown for comparison.

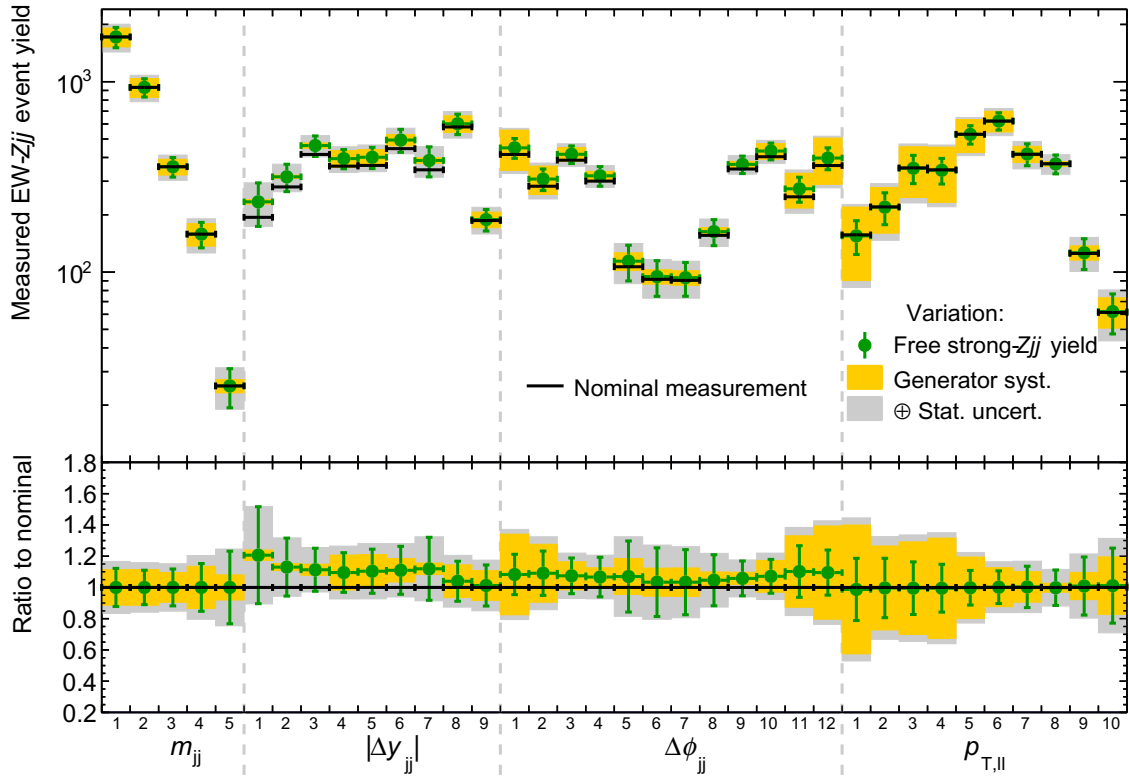


FIGURE 6.26: Measured EW- Zjj signal from the varied method without the strong- Zjj event yield constraint. The statistical uncertainty is shown as an error bar, the strong generator choice systematic is shown in orange and the combined statistical plus generator uncertainty is shown in grey. All uncertainties are derived from bootstraps and are from the varied method, the central value of the nominal result is shown for comparison.

are compared to the nominal detector-level result in Figure 6.26. Without and without the constraint the $p_T^{\ell\ell}$ measurement converges to a similar yield. The $|\Delta y_{jj}|$ and $\Delta\phi_{jj}^{\text{signed}}$ measurements prefer a lower strong- Zjj yield, which results in an increased EW- Zjj event yield. In each case the varied result is compatible with the nominal one within uncertainty.

6.6.4 Sequential alternative method

The final alternative that was studied is an independent method of performing the EW- Zjj extraction. This sequential method is a relatively simple approach where the constraints are derived and applied one-by-one to correct the strong- Zjj prediction. The purpose of this alternative is to show that an independent, more conventional approach to extract the EW signal can be used to obtain a consistent result. The same four regions are used as in the nominal method, to explain the sequential approach the quantity $X(r)$ will again be useful. Recall that in each bin i , $X(r, i)$ is defined as

$$X(r, i) = \frac{N_{r,i}^{\text{data}} - \nu_{ri}^{\text{EW,MC}} - \nu_{ri}^{\text{other,MC}}}{\nu_{r,i}^{\text{strong,MC}}}. \quad (6.11)$$

In the nominal method the correction applied to the strong- Zjj component in the SR is, from Equation (6.3),

$$\nu_{\text{SR},i}^{\text{strong}} = b_{\text{L},i} f(x_i) \nu_{\text{SR},i}^{\text{strong,MC}}. \quad (6.12)$$

Similarly, in the sequential method the correction applied is

$$\nu_{\text{SR},i}^{\text{strong}} = \beta_{r,i} k \nu_{\text{SR},i}^{\text{strong,MC}}, \quad (6.13)$$

where $\beta_{r,i}$ is the value of a function evaluated in bin i , determined from a fit to $X(r)$ in region control region r and k is an overall normalization factor.

As in the nominal method, either CRa or CRc can be used to derive the bin dependent correction that is applied to the SR. However unlike the nominal method,

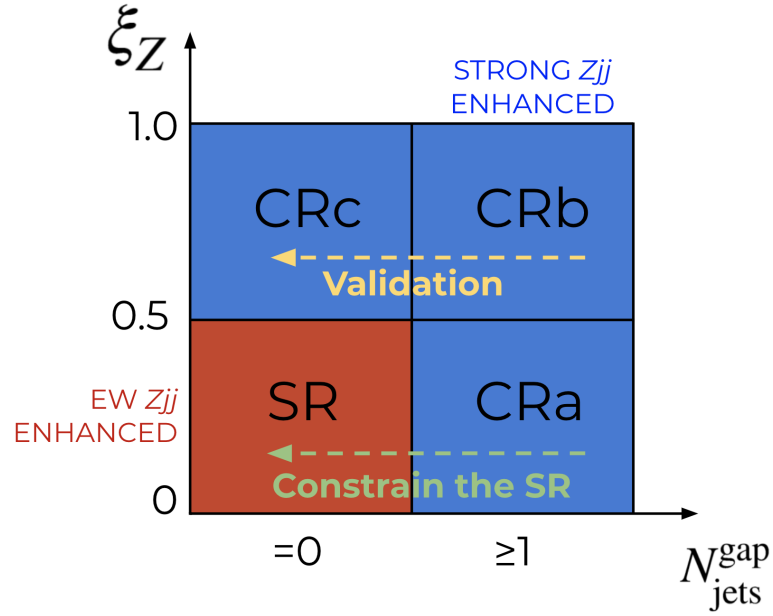


FIGURE 6.27: Diagram depicting the alternative sequential method. In this case the constraints applied to the signal region are derived in CRa and the validation is performed by applying constraints derived in CRb to CRc.

the other two control regions will not be used to derive constraints. Instead the two other control regions serve to validate the method, for example in the case that CRa is used to constrain the SR the validation will be performed by using CRb to constrain CRc as depicted in Figure 6.27. The purpose of the validation is to determine the non-closure of the sequential method. The non-closure is a result of the fact that the constraint derived in CRb will not perfectly correct the strong- Z_{jj} component in CRc, the residual data / MC ratio after correction is interpreted as an additional non-closure systematic.

The normalization factor k is determined in the normalization region, defined as a single bin in the m_{jj} spectrum defined between $500 < m_{jj} < 750$ GeV. The normalization region is essentially an additional 0th m_{jj} bin in the regions where

Variable	$\beta_r(x)$ parameterization
m_{jj}	$\beta_r(m_{jj}) = e^{a+b m_{jj}}$
$ \Delta y_{jj} $	$\beta_r(\Delta y_{jj}) = a + b \Delta y_{jj} $
$\Delta\phi_{jj}^{\text{signed}}$	$\beta_r(\Delta\phi_{jj}^{\text{signed}}) = a + b \left(\Delta\phi_{jj}^{\text{signed}}\right)^2$
$p_T^{\ell\ell}$	$\beta_r(p_T^{\ell\ell}) = a + b p_T^{\ell\ell} + c \left(p_T^{\ell\ell}\right)^2$

TABLE 6.6: Parameterization of the fit functions used in the sequential method.

the constraints are derived and applied, defined with the same $N_{\text{jets}}^{\text{gap}}$ and ξ_Z criteria as in the particular region. In the example where CRa is used to constrain the SR, the normalization factor k is defined as

$$k = \frac{X(\text{SR}, 0)}{X(\text{CRa}, 0)}, \quad (6.14)$$

the ratio of $X(r)$ in the 0th bin of the two regions.

The fit function $\beta_r(x)$ is fit to $X(r)$, the form of the function varies depending on the measurement variable. The fit function forms are summarized in Table 6.6.

As in the nominal method, multiple strong- Zjj predictions are available. The sequential method is performed twice, once with each of the SHERPA and MG5+PY8 predictions, both using the HERWIG+VBFNLO EW- Zjj prediction. The central value is the numerical average of these two results and the generator choice systematic is defined as half the difference between the values extracted with SHERPA and MG5+PY8. CRa is used to constrain the SR for each variable, except for $\Delta\phi_{jj}^{\text{signed}}$ where CRc is used. This is because the non-closure observed when using CRa with $\Delta\phi_{jj}^{\text{signed}}$ is unacceptably large. The results from the variation are compared to the nominal detector-level result in Figure 6.28. The observed agreement is reasonable,

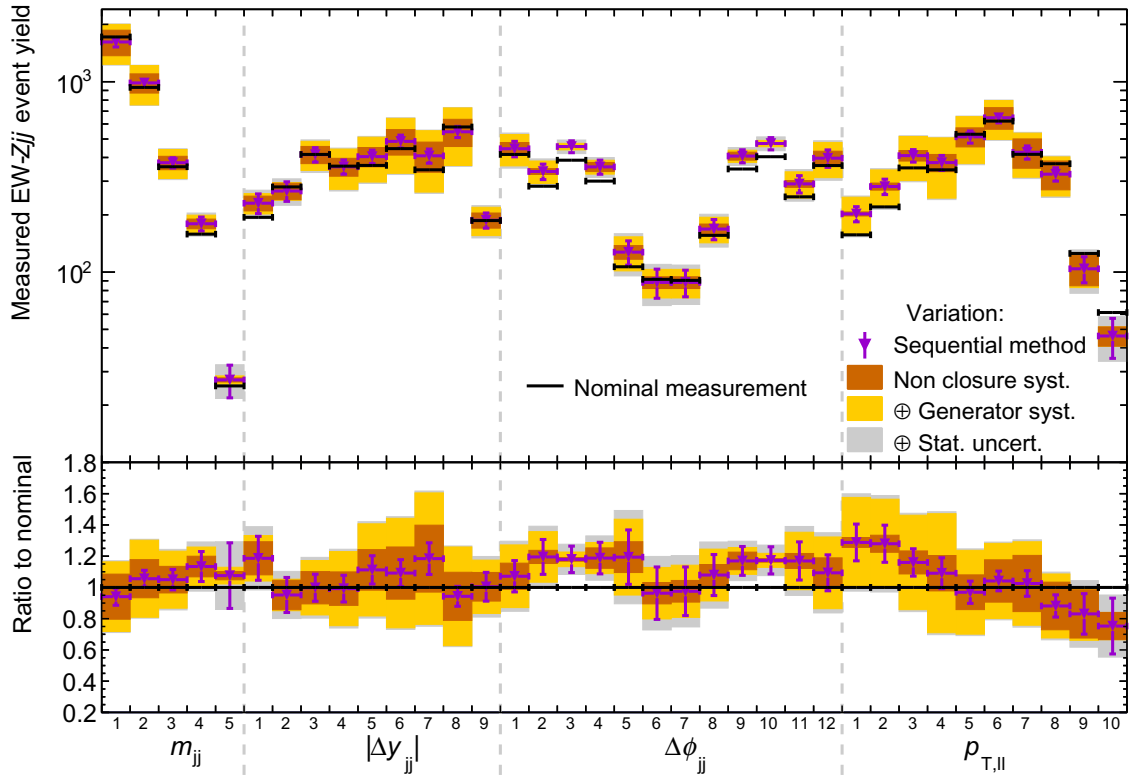


FIGURE 6.28: Measured EW- Zjj signal from the sequential method of EW- Zjj signal extraction. The statistical uncertainty is shown as an error bar, the strong generator choice systematic is shown in orange, the non-closure is shown in dark orange and the total uncertainty is shown in grey. All uncertainties are derived from bootstraps and are from the varied method, the central value of the nominal result is shown for comparison.

this will be quantified in the following section.

6.6.5 Compatibility of varied results

In this section the compatibility of the varied and nominal results will be determined. This was also achieved via bootstrapping, for each pseudo-experiment the difference in extracted EW signal per bin was computed, ΔN_{EW}^i . The mean difference from the

distribution of the pseudo-experiment results form a vector ΔN_{EW} . The covariance matrix C_{EW} is also determined from the pseudo-experiments. The difference in the generator choice systematic between the varied and nominal case is included in this covariance matrix, and for the sequential alternative the non-closure systematic is included as well. With this vector ΔN_{EW} and the covariance matrix, the χ^2 is computed by inverting the matrix as

$$\chi^2 = (\Delta N_{EW})^T C_{EW}^{-1} \Delta N_{EW}. \quad (6.15)$$

The results of these computations for each variation / alternative are presented in Table 6.7. Typical values are observed from these compatibility tests, overall the nominal method is stable to these variations and compatible results can be obtained from a wide variety of modifications. This also shows that the relatively simple sequential method can be used to produce a consistent result. Figure 6.29 shows the four varied and the nominal results together for comparison.

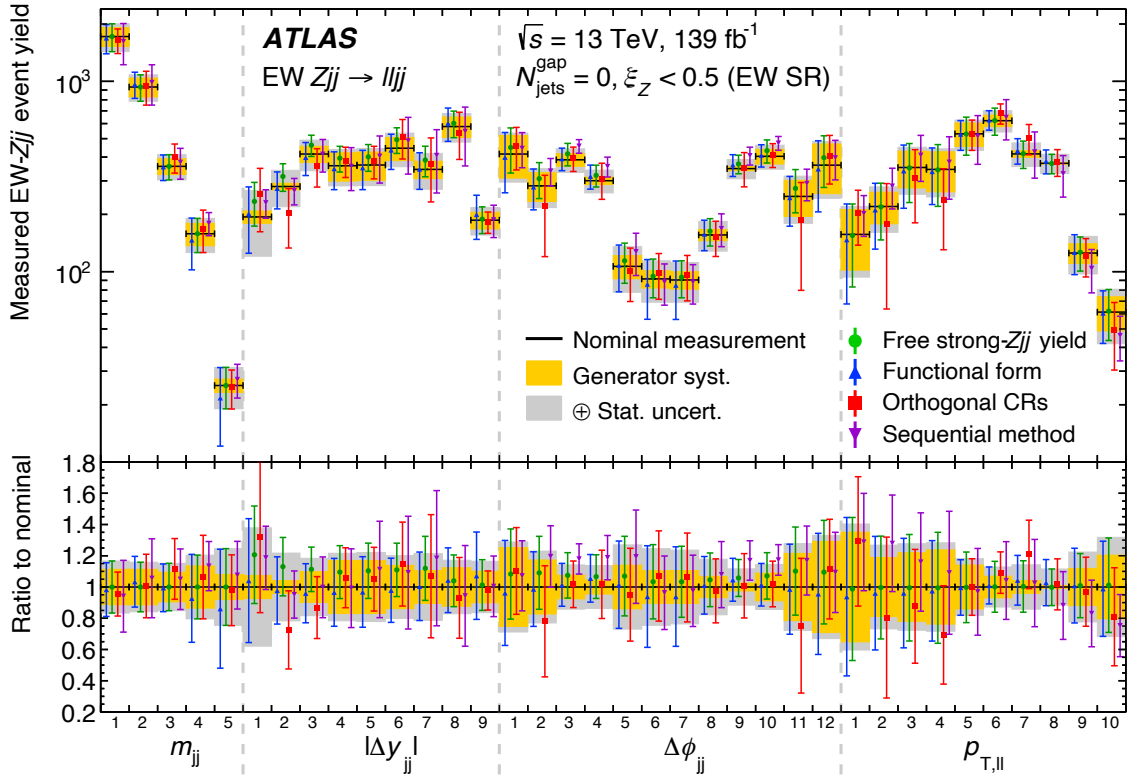


FIGURE 6.29: Measured EW- Zjj signal from the nominal method and from each variation / alternative. The error bands correspond to the nominal uncertainties, the error bars for each variation include the statistical and generator systematic. The sequential alternative also includes the non-closure systematic.

Variation / Alternative	χ^2 / d.o.f. and p -value			
	m_{jj}	$ \Delta y_{jj} $	$\Delta\phi_{jj}^{\text{signed}}$	$p_T^{\ell\ell}$
Orthogonal CRs	2.4 / 5	6.3 / 9	7.0 / 12	7.2 / 10
	0.79	0.71	0.86	0.71
Functional form	0.5 / 5	2.8 / 9	2.7 / 12	6.3 / 10
	0.99	0.97	0.99	0.79
Free strong- Zjj yield	N/A	3.9 / 9	12.8 / 12	6.9 / 10
		0.92	0.38	0.73
Sequential method	1.5 / 5	13.2 / 9	14.9 / 12	11.8 / 10
	0.91	0.15	0.25	0.30

TABLE 6.7: Results of the compatibility tests between the varied and nominal method. The free strong- Zjj yield result for the m_{jj} case is marked as N/A since this variation has no impact by construction.

6.7 Correction for detector effects

6.7.1 Iterative Bayesian unfolding

The purpose of this correction procedure is to correct the detector-level measurements to the particle-level. The detector-level measurements are unfolded, with the goal being to remove the dependence of the ATLAS detector such that they can be directly compared to theoretical predictions. MC simulations are used to model the relationship between the underlying particle-level events and the associated simulated detector-level reconstructed events. In particular, the simulation is used to correct for (i) events that migrate between bins of the measured distributions between particle and detector-level, (ii) events that pass detector-level selection but not particle-level selection and (iii) events that pass particle-level selection but not detector-level selection. The primary cause of each of these effects are (i) the finite

resolution of the detector, (ii) pileup that produces *fake* events at detector-level and (iii) inefficiency of the detector.

The iterative Bayesian unfolding procedure [110, 111] used to unfold the results will be briefly described, based on the discussion of Reference [112]. The aim of the method is to measure the differential distribution of a particle-level observable t of the signal process starting from the detector-level distribution r (*reconstructed*). The measurement is performed in N consecutive bins defined by bin boundaries $t_i^{\min}$ and $t_i^{\max}$, with $t_i^{\max} = t_{i+1}^{\min}$. Similarly, the detector-level distribution r are in N bins defined by $r_j^{\min}$ and $r_j^{\max}$, with $r_j^{\max} = r_{j+1}^{\min}$. For the measurements in this thesis, the bins span the same interval, *i.e.* $r_i^{\min} = t_i^{\min}$.

Vector $\mathbf{n}$ represents the particle-level event yields and $\mathbf{m}$ the corresponding detector-level event yields, both in N bins. These yields can be translated to a differential cross section according to Equation 5.1. Here $\mathbf{n}$ represents what a perfect detector would observe and $\mathbf{m}$ represents what is observed in the real detector (for example the EW- Zjj distribution obtained from the signal extraction procedure). The goal of the unfolding procedure is to translate the observed yields $\mathbf{m}$ into the particle-level yields $\mathbf{n}$.

There exists some matrix $\mathbf{R}$, the *response matrix*, such that for the j^{th} detector-level bin

$$m_j = \sum_i R_{ij} n_i + \beta_j, \quad (6.16)$$

where i denotes the particle-level bin. β_j denote the *out-of-fiducial* events, those that pass the analysis selection but not the particle-level fiducial selection. Thus, the response matrix transforms the underlying true yields into the observed detector

yields. It is defined as the probability that detector-level quantity r belongs to bin j given that particle-level quantity t belongs in bin i , *i.e.*

$$R_{ij} = p\left(r \in (r_j^{\min}, r_j^{\max}) | t \in (t_i^{\min}, t_i^{\max})\right). \quad (6.17)$$

For simplicity and compactness this will be written as simply

$$R_{ij} = p(r_j | t_i). \quad (6.18)$$

Using MC simulation one can determine the *migration matrix* $\mathbf{M}$, where the element M_{ij} represents the joint probability to observe the event in bin j of the reconstructed distribution and bin i of the true distribution, *i.e.*

$$M_{ij} = p\left(r \in (r_j^{\min}, r_j^{\max}), t \in (t_i^{\min}, t_i^{\max})\right) \equiv p(r_j, t_i). \quad (6.19)$$

Each event that passes the event selection at both particle-level and detector-level enters into the migration matrix. The response and migration matrices are related since the underlying probabilities are related as

$$p(r_j, t_i) = p(r_j | t_i) \cdot p(t_i). \quad (6.20)$$

This implies that

$$M_{ij} = R_{ij} \cdot p(t_i). \quad (6.21)$$

This can be combined with the reconstruction efficiency ϵ_i , the probability to reconstruct a particle-level event in bin i in any of the detector-level bins, which is given

by summing the possible outcomes,

$$\epsilon_i = \sum_j^N R_{ij} = \sum_j^N p(r_j|t_i). \quad (6.22)$$

Then summing over the i bins on both sides of Equation (6.21) gives

$$\sum_j^N M_{ij} = \epsilon_i \cdot p(t_i), \quad (6.23)$$

and substituting this for $p(t_i)$ in Equation (6.21) relates the two matrices as

$$R_{ij} = \frac{\epsilon_i \cdot M_{ij}}{\sum_k^N M_{kj}}. \quad (6.24)$$

This shows that the response matrix, which unfolding aims to invert, can be determined from the migration matrix and reconstruction efficiencies. These are both obtainable from the MC simulation.

The iterative Bayesian unfolding procedure starts from the statement of Bayes's theorem in terms of the i^{th} and j^{th} bins,

$$p(t_i|r_j) \propto p(r_i|t_j) \cdot p(t_j). \quad (6.25)$$

For experimentally measured yields $\mathbf{m}$, the expected true distribution is calculated as

$$\hat{n}_i = \sum_j^N m_j \cdot p(t_i|r_j), \quad (6.26)$$

and taking into account the inefficiency

$$\hat{n}_i = \frac{1}{\epsilon_i} \sum_j^N m_j \cdot p(t_i|r_j). \quad (6.27)$$

In the iterative approach the conditional probabilities $p(t_i|r_j)$ form the *unfolding matrix*, $\mathbf{U}$. This matrix is built from $\mathbf{M}$ and from a prior probability distribution. The initial prior is taken to be the nominal MC prediction of the signal process, in this case EW-Zjj. From Bayes's theorem

$$U_{ij} = p(t_i|r_j) = \frac{p(r_j|t_i) \cdot p(t_i)}{\sum_k^N p(r_j|t_k) \cdot p(t_k)}, \quad (6.28)$$

here the response matrix $p(r_j|t_i)$ is taken from the migration matrix $\mathbf{M}$ determined in MC and the $p(t_i)$ are the signal MC prior. The unfolding matrix is then applied to the observed detector-level distribution to obtain the first unfolded particle-level distribution

$$\hat{n}_i = \frac{1}{\epsilon_i} (\mathbf{U} \cdot \mathbf{m})_i. \quad (6.29)$$

This is iterated by rederiving the unfolding matrix using the result obtained $\hat{\mathbf{n}}$ as the updated prior. The more this is iterated, the less the result will be bias towards the starting MC prediction. However, there is a trade off, as the bias decreases the statistical uncertainty increases with subsequent applications of the migration matrix. Thus in general the number of iterations will be small (< 10) and it was found that two iterations is sufficient to unfold the results of this analysis. The first closure test, described in the next section, will show that the unfolded result is relatively insensitive to the particular MC generator used by switching to an

alternative generator.

6.7.2 Implementation

With the general unfolding method outlined, this section will define the distributions and quantities used for unfolding in detail. As previously defined, the event yields $\mathbf{n}$ at particle-level and $\mathbf{m}$ at detector-level, these will each have N bins as defined for the measurement variables in Table 5.6. Here $\mathbf{m}$ is the observed event yields, either the measured EW- Zjj yields from the signal extraction or data minus non- Zjj background for the inclusive Zjj measurement. The quantities used in unfolding, which are determined from the MC simulation are:

- The migration matrix $\mathbf{M}$, is used to correct for bin-to-bin migrations, the matrix is filled for each event that passes the event selection at detector *and* particle levels.
- The fiducial factors $\mathcal{F}_j = \frac{m_j - \beta_j}{m_j}$ defined for each bin j determined by the value at detector-level. This is the ratio of the number particle-level events that fall within the fiducial definition to the number of associated detector-level events in bin j . This factor corrects for the fake reconstructed events that do not fall into the fiducial definition at particle-level, $0 \leq \mathcal{F}_i \leq 1$.
- The efficiency correction factors ϵ_i defined for each bin i determined by the value at particle-level. This is the ratio of the number of reconstructed events that are correctly reconstructed to the true number of associated particle-level events in bin i as in Equation 6.23. This factor corrects for the detector reconstruction efficiency, $0 \leq \epsilon_i \leq 1$.

Based on these distributions and quantities the signal and unfolded distributions are defined as

$$\begin{aligned} m'_j &= m_j \cdot \mathcal{F}_j, \\ n_i &= \frac{1}{\epsilon_i} (\mathbf{U} \cdot \mathbf{m}')_i, \end{aligned} \tag{6.30}$$

where $\mathbf{U}$ is defined from $\mathbf{M}$ and the prior as described previously. The default strong- Zjj sample used for unfolding is the one from the SHERPA generator and the default EW- Zjj sample is from POWHEG+PY8. Examples of the migration matrices are shown in Figure 6.30, these are from the unfolding of the inclusive Zjj spectrum in CRc. The fiducial factors $\mathcal{F}_j$ are typically in the range of 0.7 to 0.9, and the efficiency correction factors are typically between 0.4 and 0.6. Figure 6.31 shows the unfolded result after various numbers of iterations ranging from two to ten. The results tend to converge rapidly and increasing the number of iterations beyond two was not found to be worthwhile considering the increased statistical uncertainty resulting from subsequent applications of the migration matrix.

6.7.3 Closure tests and systematic uncertainties

This section will outline two closure tests performed on the unfolding procedure. An unfolding systematic uncertainty component is extracted from the non-closure observed in each test.

The first closure test is designed to test the dependence of the unfolding on the particular MC simulations used in the unfolding procedure. For the inclusive Zjj measurement it is performed by repeating the procedure with the MG5+PY8 sample in place of SHERPA for the strong- Zjj component, however the MG5+PY8

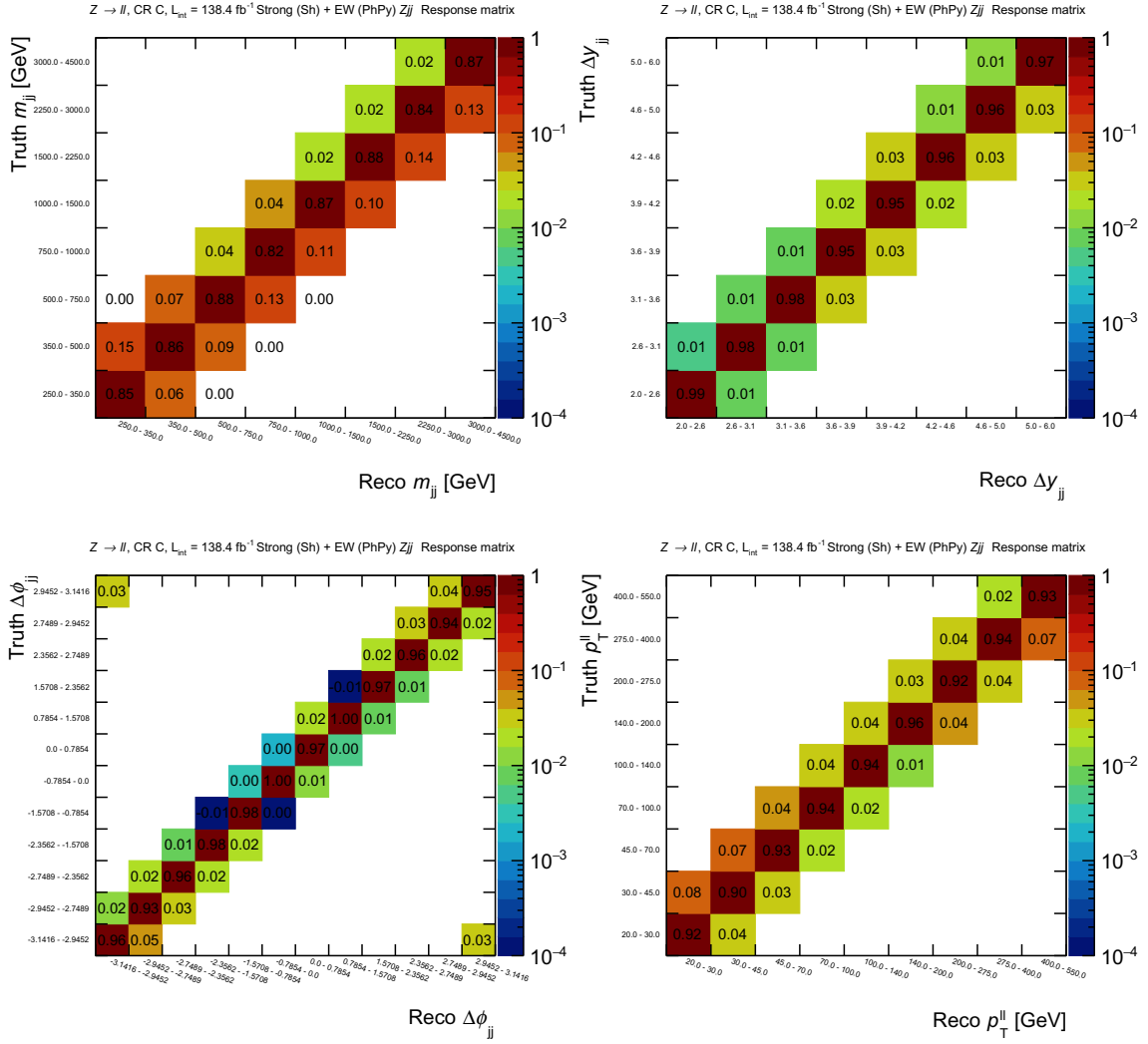


FIGURE 6.30: Migration matrices derived from the combination of SHERPA strong- Zjj and POWHEG+PY8 EW- Zjj . These matrices were used in the unfolding of the inclusive Zjj process in CRc for each of the four measurement variables.

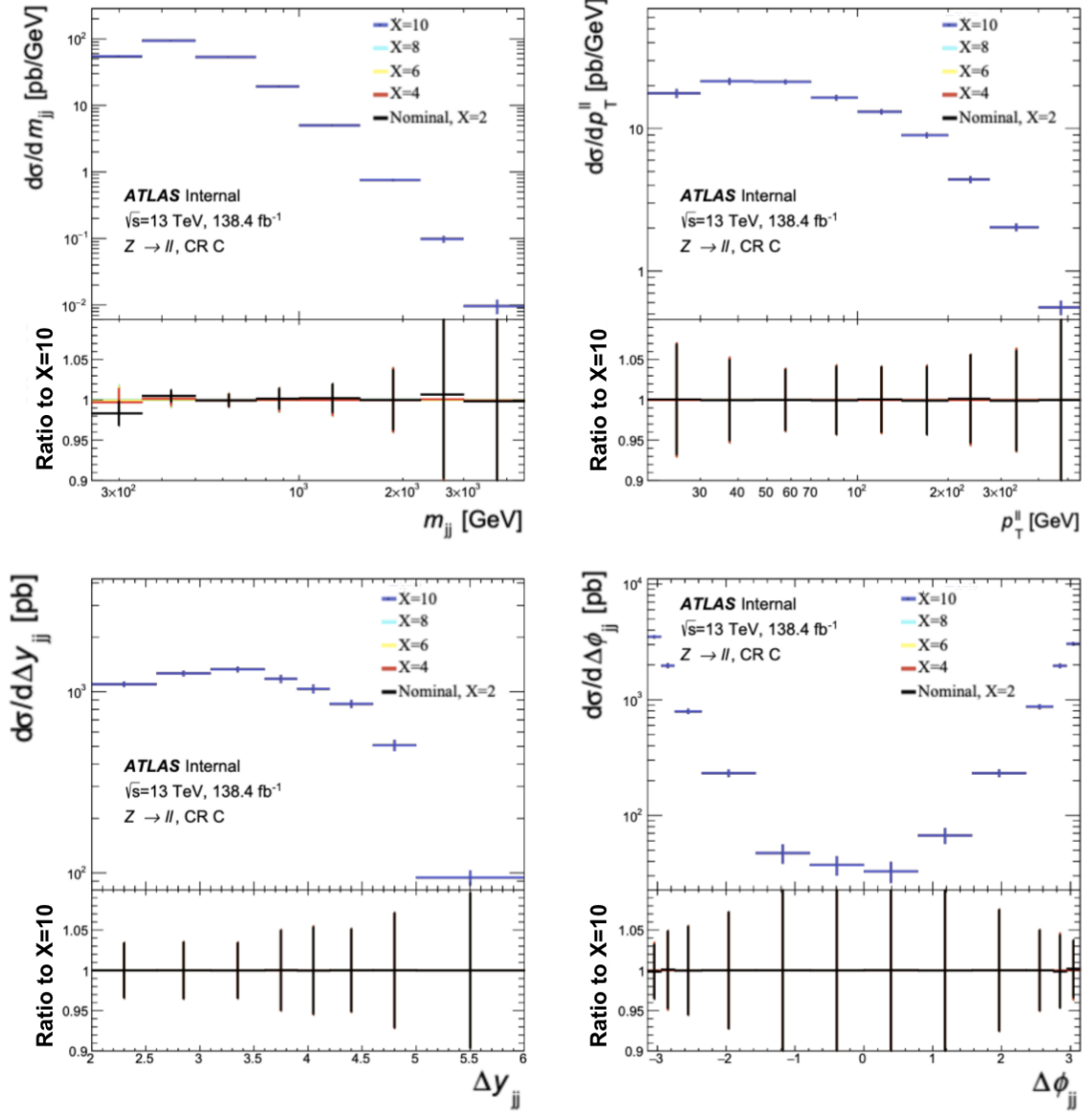


FIGURE 6.31: Impact of the number of iterations of the Bayesian unfolding on the unfolded result. Increasing the number of iterations beyond two does not significantly impact the result.

sample is reweighted such that the distribution to be unfolded matches the SHERPA distribution at particle-level. Thus the desired true distribution remains unchanged but the initial prior and unfolding matrix are modified. Similarly for the EW- Zjj process the procedure is repeated with SHERPA replacing POWHEG+PY8 for the signal prediction, with SHERPA reweighted to match POWHEG+PY8 at particle-level. A smooth reweighting is derived from the ratio of the distribution in the two samples, an example of this will be shown for the second closure test. The purpose of the reweighting is such that the initial prior distribution is essentially unchanged, therefore any difference observed in the result is due to the correction factors and migration matrix determined from the alternative generator.

The ratio of the unfolded result obtained with the alternate generator is compared to the reweighted particle-level prediction. Any difference between them is due to non-closure of the unfolding procedure and is taken as an unfolding systematic uncertainty. The result of this test from the inclusive Zjj measurement unfolding in CRa is shown in Figure 6.32. This figure also shows the smoothing applied to obtain the systematic uncertainty.

The second closure test is the *data driven* closure, this is performed to estimate the bias of the unfolding method when performed on simulated distributions similar to the real data. The simulation is reweighted at particle-level such that it provides a better description of the data at detector-level. The reweighting functions used for the data driven closure of the EW- Zjj measurement are shown in Figure 6.33, these are derived from the ratio of data to the detector-level MC prediction, the ratio is smoothed as shown in the figure.

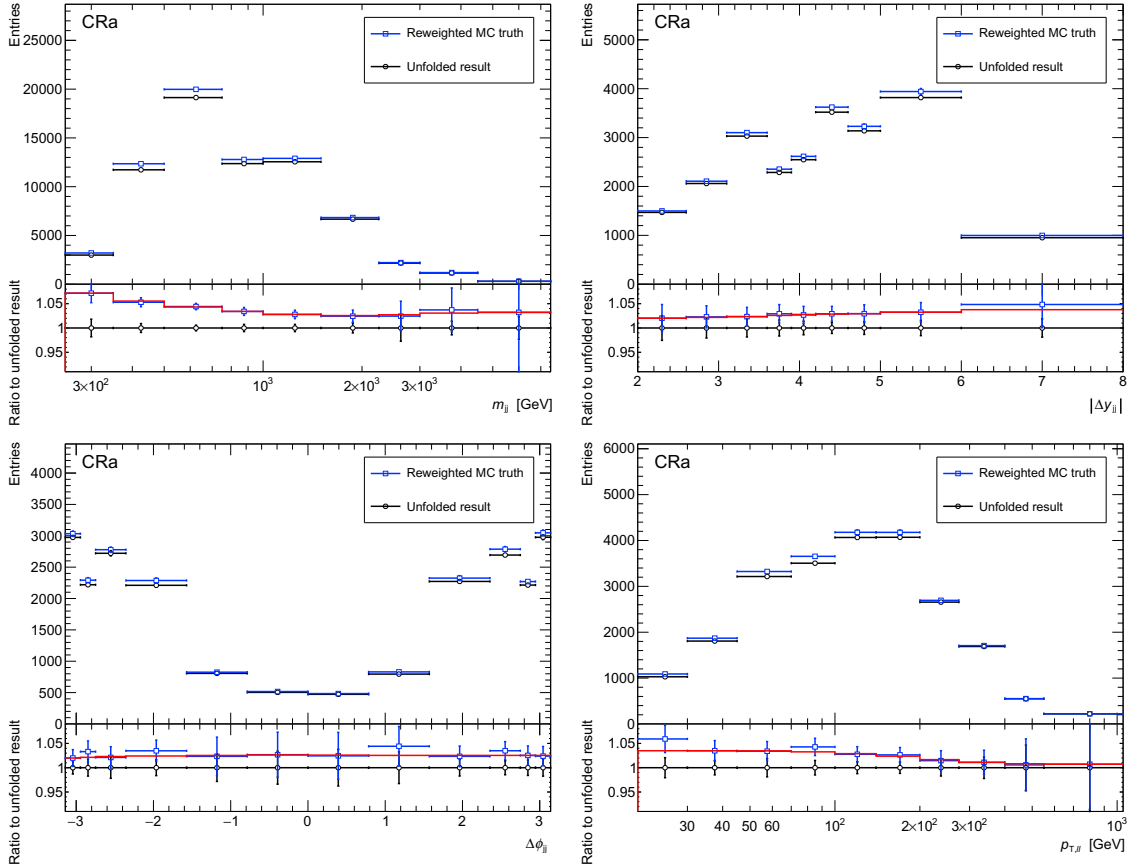


FIGURE 6.32: Result of the MC generator closure test performed for the inclusive Zjj measurement unfolding in CRa for each measurement variable. The unfolded result is compared to the reweighted particle-level prediction (truth), the ratio in the lower panel is taken as the systematic uncertainty associated to this non-closure. The final smoothed systematic uncertainty is shown in red.

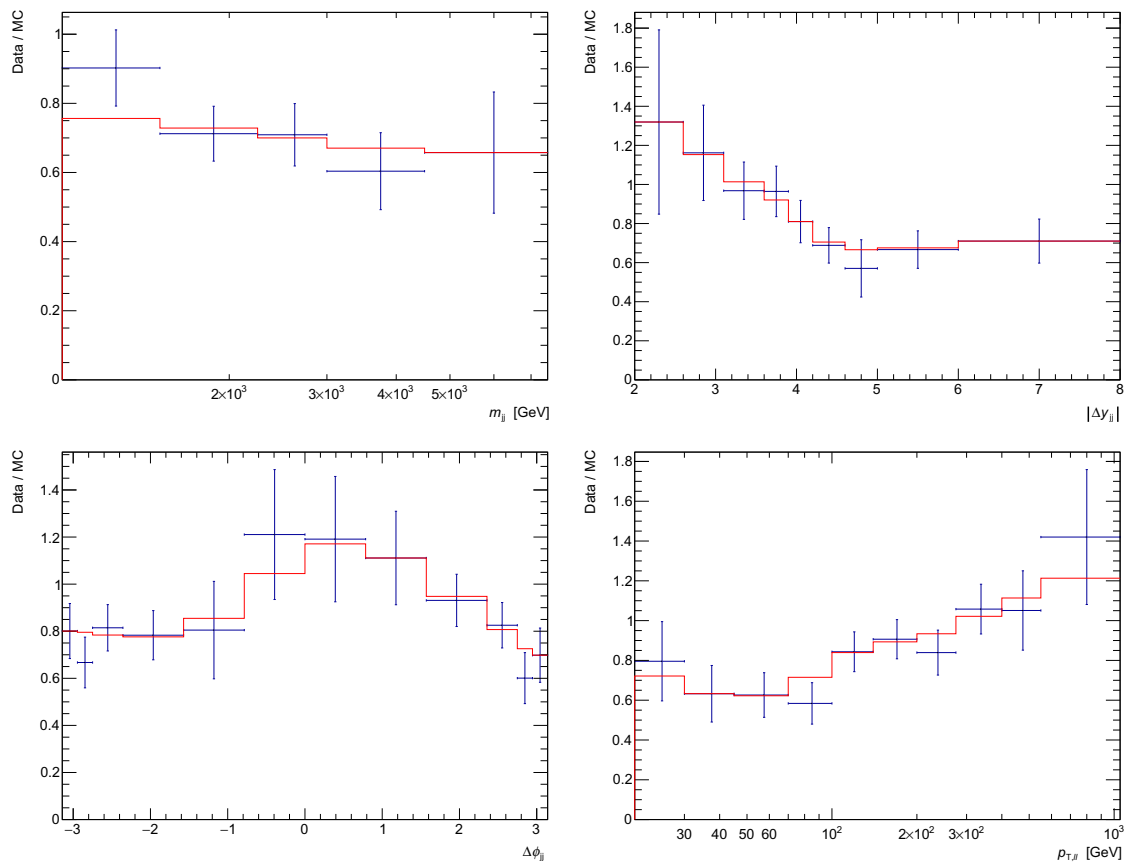


FIGURE 6.33: Ratio of the data to the detector-level MC prediction of EW-Z jj from POWHEG+PY8 in the SR for each measurement variable. The smoothed ratio (red) defines the reweighting applied to the MC to describe the data.

The reweighted detector-level prediction is then unfolded using the migration matrix and other corrections derived from nominal (unweighted) POWHEG+PY8 EW- Zjj simulation. The systematic uncertainty associated with the unfolding method is defined as the difference between the unfolded spectrum and the reweighted particle-level prediction. The result of the data driven closure for the EW- Zjj measurement is shown in Figure 6.34, as before the final systematic uncertainty has been smoothed.

Besides the systematic uncertainties associated to the unfolding procedure itself, the experimental and theoretical uncertainties described in Section 6.4 are propagated through the unfolding procedure. Each source is separately propagated by repeating the unfolding procedure after varying the MC simulation, the change in the unfolded event yield compared to the nominal result is taken as the impact of the systematic uncertainty. The statistical uncertainty on the unfolded result is determined with the extracted EW- Zjj signal results from the pseudo-experiments described in Section 6.4.3. Each pseudo-experiment is unfolded separately and the RMS spread of the distribution of the unfolded pseudo-experiment results is taken as the statistical uncertainty in each bin.

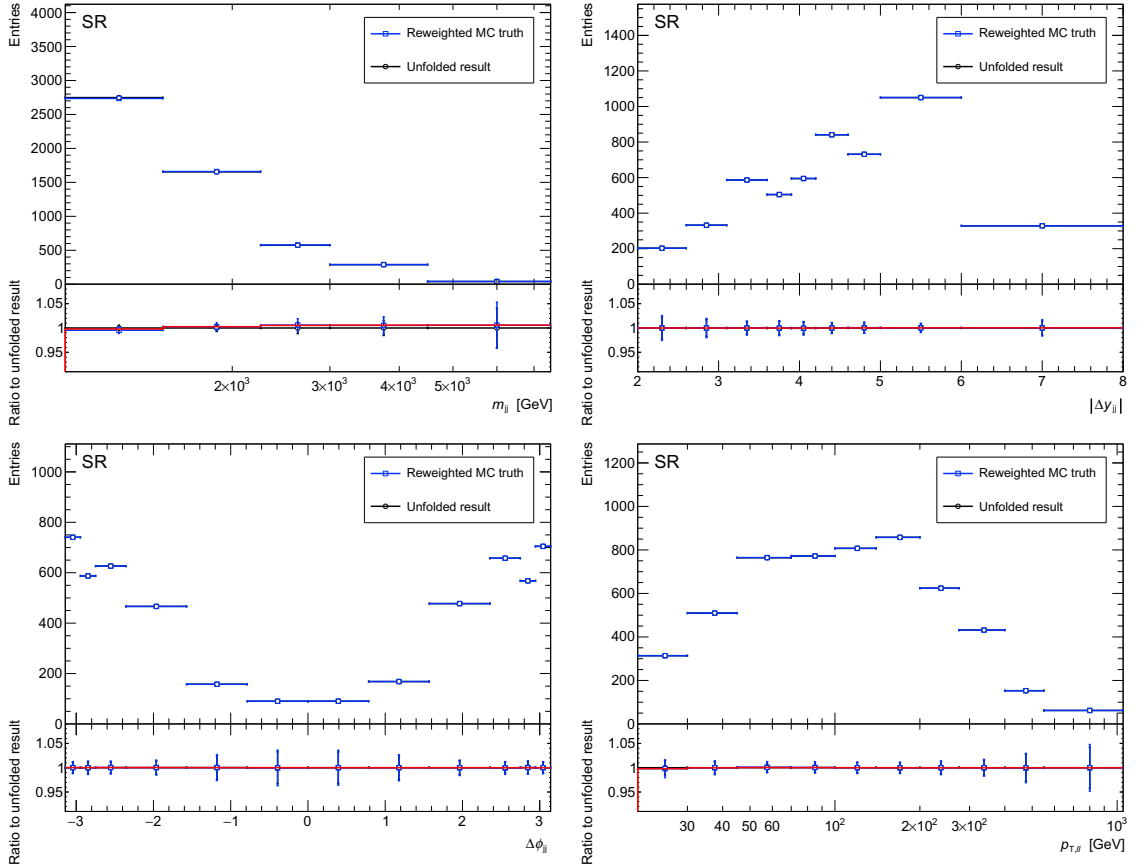


FIGURE 6.34: Result of the data driven closure performed for the EW-Zjj measurement unfolding in SR for each measurement variable. The unfolded result is compared to the reweighted particle-level prediction (truth), the ratio in the lower panel is taken as the systematic uncertainty associated to this non-closure. The final smoothed systematic uncertainty is shown in red.

6.8 Unfolded results

The final unfolded differential cross sections for EW- Zjj production as a function of m_{jj} , $|\Delta y_{jj}|$, $\Delta\phi_{jj}^{\text{signed}}$ and $p_T^{\ell\ell}$ are presented in Figure 6.35. These figures compare the results to the particle-level predictions from SHERPA, POWHEG+PY8 and HERWIG+VBFNLO. The final uncertainties of the differential cross section measurements of EW- Zjj are shown in Figure 6.36, the uncertainties associated with the modelling of the strong- Zjj process dominate.

The HERWIG+VBFNLO prediction is found to be in better agreement with the data compared to the other EW- Zjj generators. The POWHEG+PY8 prediction overestimates the data at high m_{jj} , high $|\Delta y_{jj}|$ and intermediate $p_T^{\ell\ell}$. The POWHEG+PY8 prediction often does not agree with the HERWIG+VBFNLO prediction within the theoretical uncertainty. A similar discrepancy between theoretical predictions was noted for EW $VVjj$ processes in Ref. [74] and was attributed to the set-up of the parton shower when matched to the matrix-element calculations. The SHERPA prediction significantly underestimates the measured differential cross-sections, due to a non-optimal setting of the colour flow [74].

The total fiducial cross sections are calculated for each measurement variable by integrating the differential cross-section. The fiducial cross section measured from the combination of the four measurement observables is 37.14 ± 3.47 (stat.) ± 5.79 (syst.) fb. Table 6.8 presents the measured cross sections along with the p-value calculated for the agreement to the theoretical prediction from HERWIG+VBFNLO, which is 39.5 ± 3.4 (scale) ± 1.3 (PDF) fb. The agreement with the HERWIG+VBFNLO prediction is found to be good for each variable.

Also presented in Table 6.8 is the fiducial cross section measured from the

combination of the four measurement variables. The uncertainty on the combined cross section is determined via the full correlation matrices presented in Figure 6.37 (statistical only) and Figure 6.38 (stat. plus syst.). These correlation matrix include the cross correlations between the bins of different measurement variables. The statistical and strong- Zjj generator choice systematic correlations are determined via the unfolded pseudo-experiment results.

The final unfolded differential cross sections for inclusive Zjj production as a function of m_{jj} , $|\Delta y_{jj}|$, $\Delta\phi_{jj}^{\text{signed}}$ and $p_T^{\ell\ell}$ are presented in Figure 6.39, also in the EW signal region. These figures compare the results to the particle-level predictions from SHERPA and MG5_NLO+PY8. The final uncertainties of the inclusive measurement in the SR are shown in Figure 6.40, the jet energy scale and jet energy resolution uncertainties dominate. Additional differential cross section and uncertainty figures for the inclusive measurements performed in the three control regions are presented in Appendix B.

The unfolded differential cross sections for both EW- Zjj production and inclusive Zjj production are documented in tabular form, also in Appendix B. The data are also provided in a HEPDATA [112] repository [113] with a fully granular breakdown of the systematic uncertainty sources.

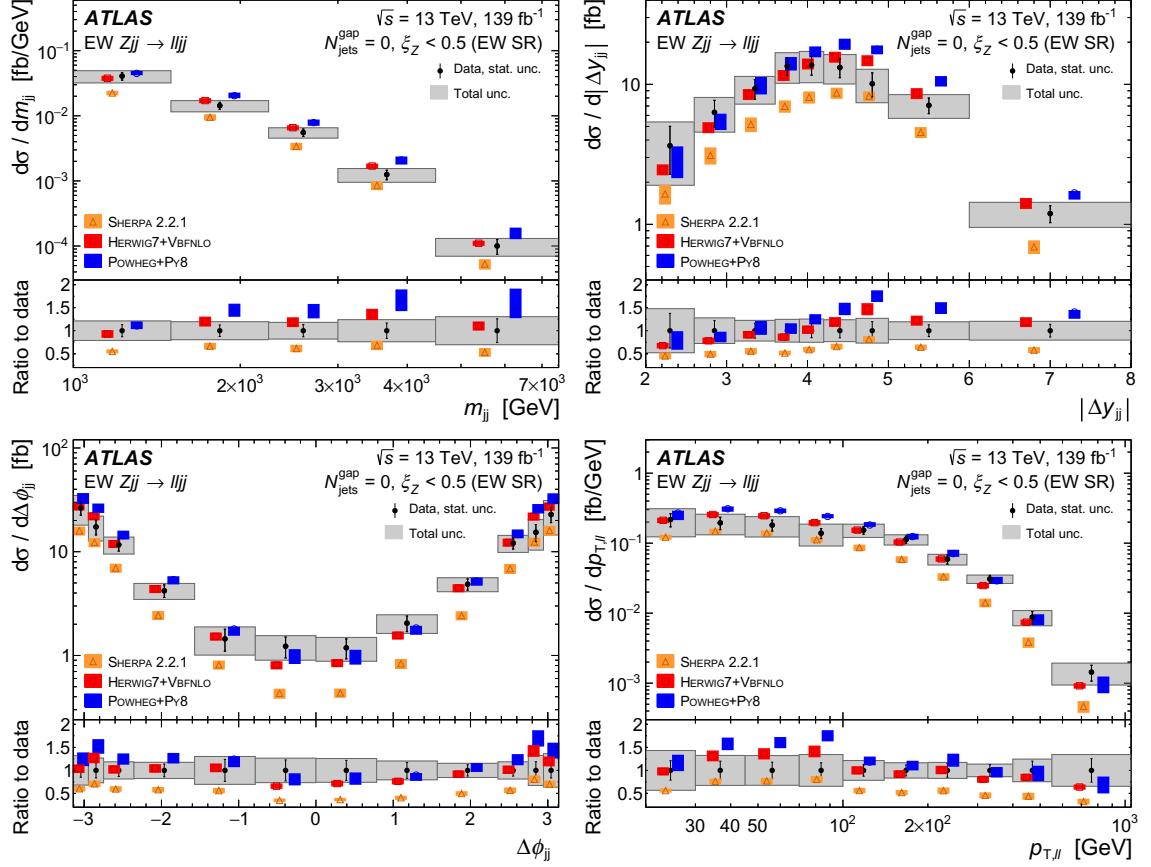


FIGURE 6.35: Differential cross sections for EW-Z jj production as a function of the four measurement variables. The unfolded data are shown as black points, with the statistical uncertainty represented by an error bar and the total uncertainty represented as a grey band. The data are compared with theoretical predictions produced by HERWIG + VBFNLO (red points), POWHEG+PY8 (blue points) and SHERPA (orange points). Uncertainty bands are shown for the three theoretical predictions. Each theory prediction is slightly offset from the bin centre to avoid overlap.

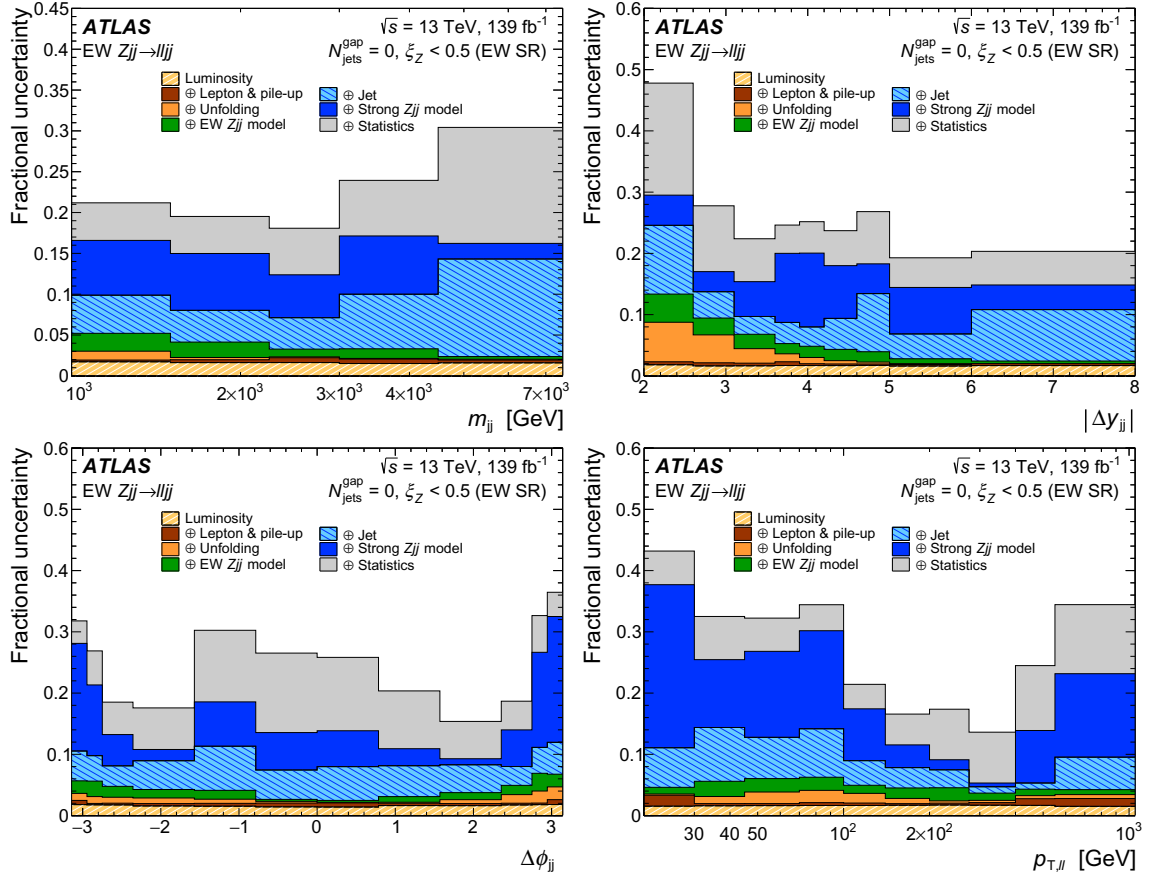


FIGURE 6.36: Fractional uncertainty in the EW-Z jj measurement as a function of the four measurement variables. Uncertainty sources are grouped in categories that are added in quadrature (denoted $\oplus$) to give the total uncertainty.

Variable	Fiducial cross section [fb]	p-value Herwig EW-Zjj
m_{jj}	37.41 ± 3.47 (stat.) ± 5.47 (syst.)	0.639
$ \Delta y_{jj} $	36.94 ± 3.44 (stat.) ± 5.81 (syst.)	0.651
$\Delta\phi_{jj}^{\text{signed}}$	37.21 ± 3.50 (stat.) ± 5.88 (syst.)	0.343
$p_T^{\ell\ell}$	37.01 ± 3.48 (stat.) ± 5.85 (syst.)	0.240
Combination	37.14 ± 3.47 (stat.) ± 5.79 (syst.)	0.494

TABLE 6.8: The observed total fiducial cross sections of EW-Zjj for each measurement variable and the p-values calculated from the χ^2 for N_{bins} d.o.f. comparing the measurement to the Herwig EW-Zjj prediction. Also presented is the fiducial cross section measured from the combined result from all four observables. The HERWIG+VBFNLO predicted cross section is 39.5 ± 3.4 (scale) ± 1.3 (PDF) fb.

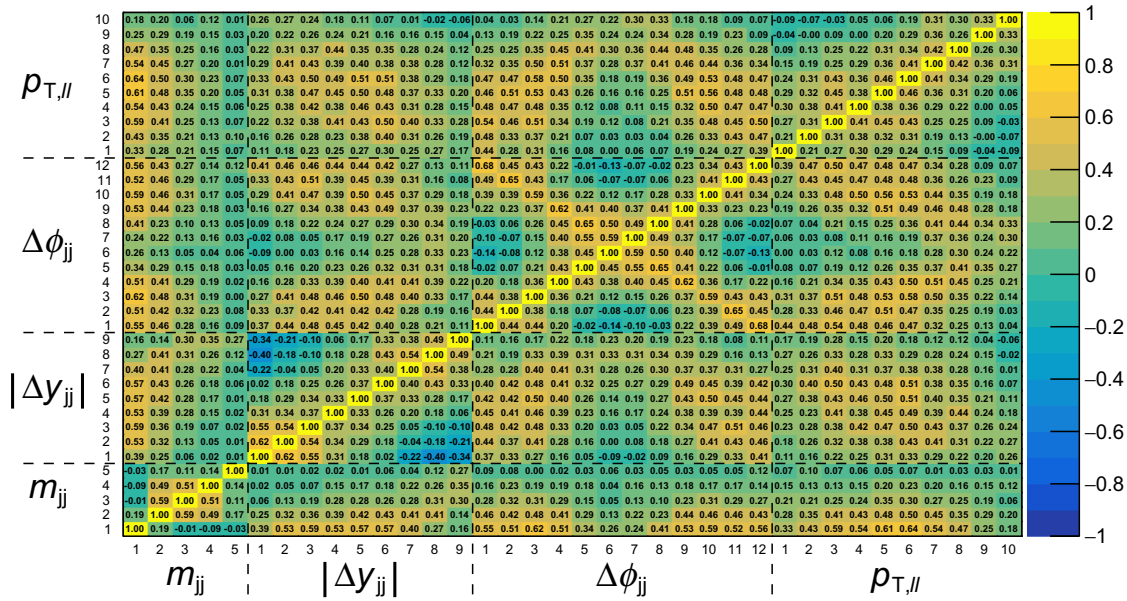


FIGURE 6.37: Statistical uncertainty correlation matrix between each bin of each measurement variable. This has been evaluated via the bootstrapping procedure outline in Section 6.4.3. The axis labels specify the bin number.

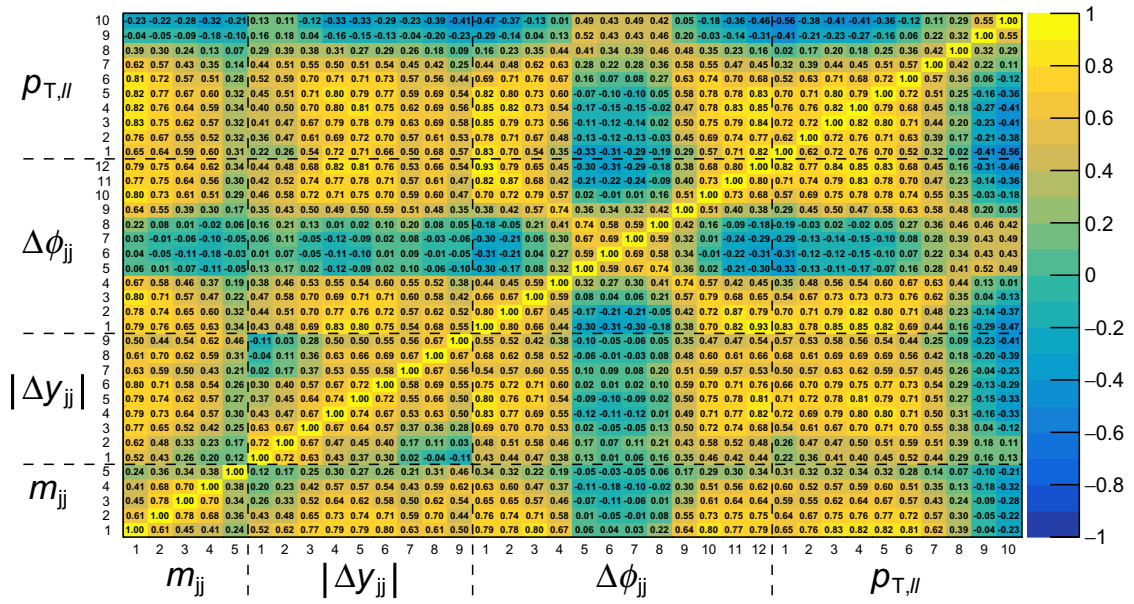


FIGURE 6.38: Full statistical plus systematic uncertainty correlation matrix between each bin of each measurement variable. This has been evaluated via the bootstrapping procedure outline in Section 6.4.3. The axis labels specify the bin number.

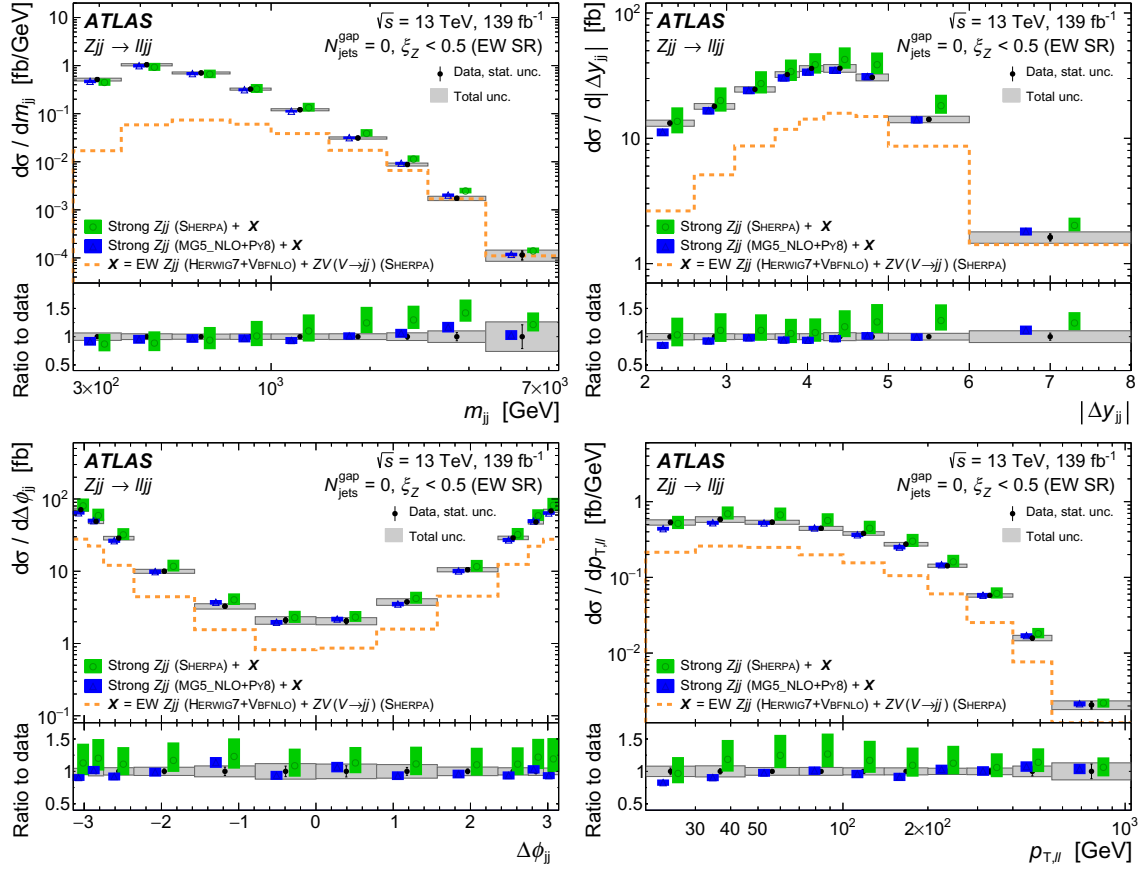


FIGURE 6.39: Differential cross-sections measured in the SR for inclusive Zjj production as a function of the four measurement variables in the signal region. The unfolded data are shown as black points, with the statistical uncertainty represented by an error bar and the total uncertainty represented as a grey band. The data are compared with theoretical predictions constructed from different strong- Zjj predictions provided by SHERPA (green) and MG5_NLO+PY8 (blue). Uncertainty bands are shown for the two theoretical predictions. Each theory prediction is slightly offset from the bin centre to avoid overlap.

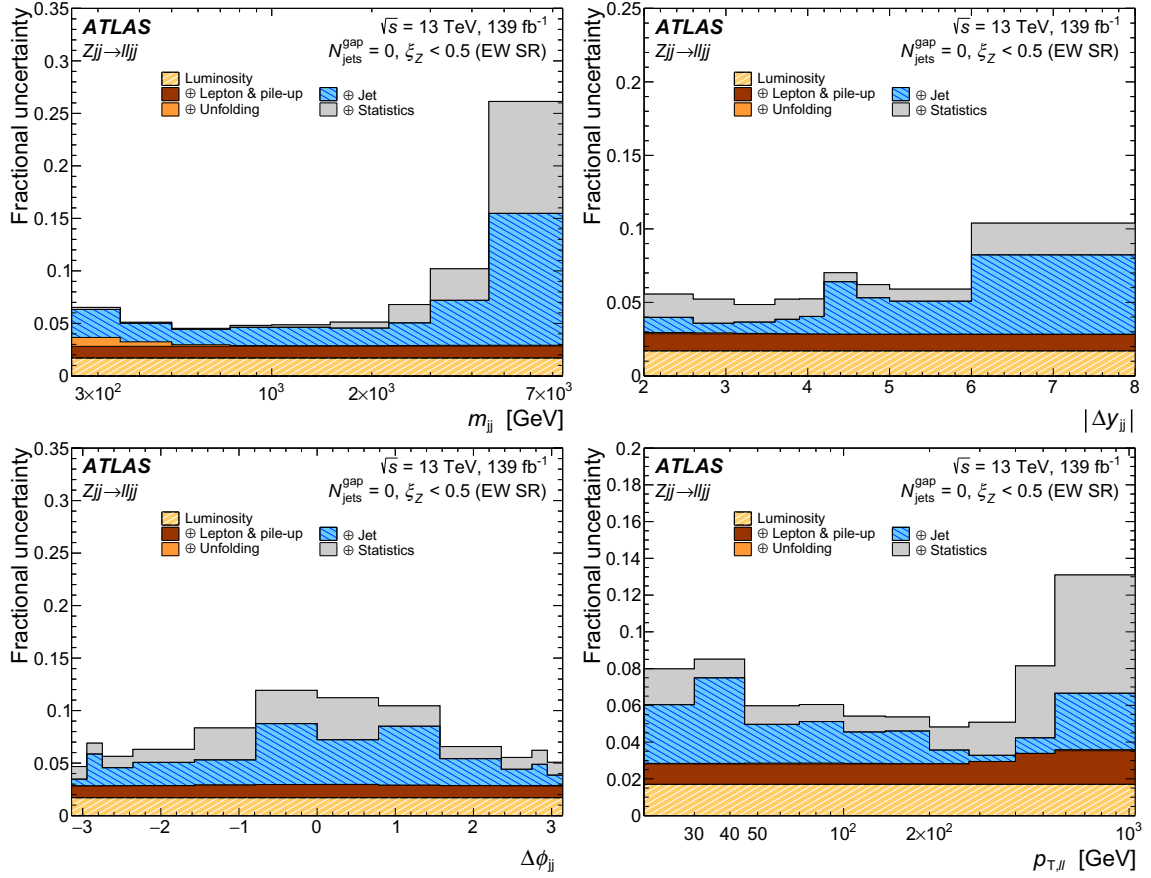


FIGURE 6.40: Fractional uncertainty in the inclusive Zjj measurement as a function of the four measurement variables in the signal region. Uncertainty sources are grouped in categories that are added in quadrature (denoted $\oplus$) to give the total uncertainty.

Chapter 7

Interpretation of results from an effective field theory perspective

The measured differential cross sections of the EW- Zjj process can be used to investigate anomalous weak-boson self interactions in the WWZ vertex. The anomalous interactions are introduced via an effective field theory as described in Section 2.6. To recap, there are four dimension six operators denoted $\mathcal{O}_W$, $\tilde{\mathcal{O}}_W$, $\mathcal{O}_{HWB}$ and $\tilde{\mathcal{O}}_{HWB}$ in the Warsaw basis that modify the WWZ vertex. The dimension six operators will impact the scattering amplitude $|\mathcal{M}|^2$ (and therefore the cross section) both with a pure dimension six contribution and via SM - dimension-six interference,

$$|\mathcal{M}|^2 = |\mathcal{M}_{\text{SM}}|^2 + 2\text{Re}(\mathcal{M}_{\text{SM}}^* \mathcal{M}_{\text{d6}}) + |\mathcal{M}_{\text{d6}}|^2. \quad (7.1)$$

The strength of the anomalous interactions are characterized by the Wilson coefficients c_k , each associated to a particular operator Q_k .

This chapter will describe the technical details of the generation of the EFT predictions and the implementation of the statistical method that extracts the best

fit values of the Wilson coefficients. The limits placed on the Wilson coefficients as a result of this study are also presented.

7.1 Generation and prediction

The pure-SM contribution to the differential cross section is taken to be the HERWIG + VBFNLO prediction, since it was found to be in good agreement with the unfolded result. The contribution from pure dimension six terms and from the SM - dimension six interference are generated at LO in perturbative QCD using MADGRAPH5 + PY8 using the NNPDF3.0NLO PDF set [73]. The parton shower, hadronization and underlying event were simulated using the A14 set of tuned parameters [81]. The interactions from the dimension six operators is provide by the SMEFTSim package [114]. Events are then passed through the Rivet routine to obtain cross-section predictions in all measured bins. Loose generator selection criteria were applied to improve generator efficiency, these were $p_T^{j1} > 65 \text{ GeV}$, $p_T^{j2} > 60 \text{ GeV}$, $m_{jj} > 200 \text{ GeV}$ and $70 < m_{\ell\ell} < 115 \text{ GeV}$.

The expected cross section is parameterized in terms of contributions from the SM (σ_{SM}), SM - dimension six interference (σ_{Linear}) that scales linearly with c_k and pure dimension six ($\sigma_{\text{Quad.}}$) that scales quadratically with c_k . If only a single Wilson coefficient c is varied at a time, the cross section will vary according to

$$\sigma_{\text{EFT}}(c) = \sigma_{\text{SM}} + \sigma_{\text{Linear}}(c) + \sigma_{\text{Quad.}}(c). \quad (7.2)$$

By evaluating σ_{Linear} and $\sigma_{\text{Quad.}}$ via simulation at a single arbitrarily-chosen value $c = c'$, the value of $\sigma_{\text{EFT}}(c)$ can be predicted for any value of c by extrapolating

according to

$$\sigma_{\text{EFT}}(c) = \sigma_{\text{SM}} + \left(\frac{c}{c'}\right) \sigma_{\text{Linear}}(c) + \left(\frac{c}{c'}\right)^2 \sigma_{\text{Quad.}}(c). \quad (7.3)$$

This decomposition method allows for a large sample of MC events to be generated at the chosen value $c = c'$ in order to benefit from high statistical-precision. The method was validated by comparing the extrapolated cross section to one determined from dedicated samples generated over a range of Wilson coefficient values. Figure 7.1 shows the expected EFT contributions to the differential cross sections using representative values of $c_W = \tilde{c}_W = 0.2$ and $c_{HWB} = \tilde{c}_{HWB} = 1.8$. The following features are of note:

- $\mathcal{O}_W[c_W = 0.2]$: For this CP-even operator, the high $p_T^{\ell\ell}$ region is particularly sensitive to the pure dimension six contribution. This feature was seen in previous studies of EW Vjj production [8, 105]. The $\Delta\phi_{jj}^{\text{signed}}$ is also very sensitive to the interference between the SM and this operator. The linear contribution is much smaller in the non- $\Delta\phi_{jj}^{\text{signed}}$ observables because each bin of these observables is integrated over the full range of $\Delta\phi_{jj}^{\text{signed}}$, which includes roughly equal positive and negative contributions.
- $\tilde{\mathcal{O}}_W[\tilde{c}_W = 0.2]$: For the CP-odd operator the impact is similar, however the linear contribution is now exactly antisymmetric about $\Delta\phi_{jj}^{\text{signed}} = 0$. The quadratic contribution is identical to that of $\mathcal{O}_W$ since it is determined by the squared pure dimension six amplitude. The linear term disappears in non- $\Delta\phi_{jj}^{\text{signed}}$ observables since the $\Delta\phi_{jj}^{\text{signed}}$ distribution is exactly symmetric and integrated over.

- $\mathcal{O}_{HWB}[c_{HWB} = 1.8]$: Another CP-even operator and the $\Delta\phi_{jj}^{\text{signed}}$ spectrum is again symmetric. There is a also significant flat linear and smaller quadratic contribution in the non- $\Delta\phi_{jj}^{\text{signed}}$ observables. This flat impact in m_{jj} and $|\Delta y_{jj}|$ is because the operator does not couple to the momentum of the W bosons. It does however couple to the Z boson momentum causing some shape in the impact at high $p_T^{\ell\ell}$.
- $\tilde{\mathcal{O}}_{HWB}[\tilde{c}_{HWB} = 1.8]$: For the CP-odd observable, as expected the $\Delta\phi_{jj}^{\text{signed}}$ spectrum is again antisymmetric about $\Delta\phi_{jj}^{\text{signed}} = 0$ and the linear contribution in the non- $\Delta\phi_{jj}^{\text{signed}}$ observables integrates to zero as before. There is no small quadratic component as in the corresponding CP-even operator.

From these observations it was concluded that the $\Delta\phi_{jj}^{\text{signed}}$ observable is the best candidate for setting limits on the EFT coefficients. This is because it is the most sensitive to each operator and the sensitivity is driven by the linear component. This is advantageous because the linear dimension six interference component is only suppressed by a factor of Λ^{-2} . The quadratic component is further suppressed as it scales as Λ^{-4} , so it and even higher dimension EFT contributions are unlikely to significantly impact the limits based primarily on the linear component.

Before extracting the EFT limits, two corrections were applied to the MADGRAPH5+PY8 EFT templates. The first is an overall normalization correction to account for the scale difference between the SM predictions of HERWIG+VBFNLO and MADGRAPH5+PY8. As shown in Figure 7.2, combining the cross section in all $\Delta\phi_{jj}^{\text{signed}}$ bins, the HERWIG+VBFNLO / MADGRAPH5+PY8 ratio is 1.104. Thus, a normalization k factor of +10.4% is applied to the EFT templates from the LO MADGRAPH predictions to account for higher order QCD effects. The figure also shows

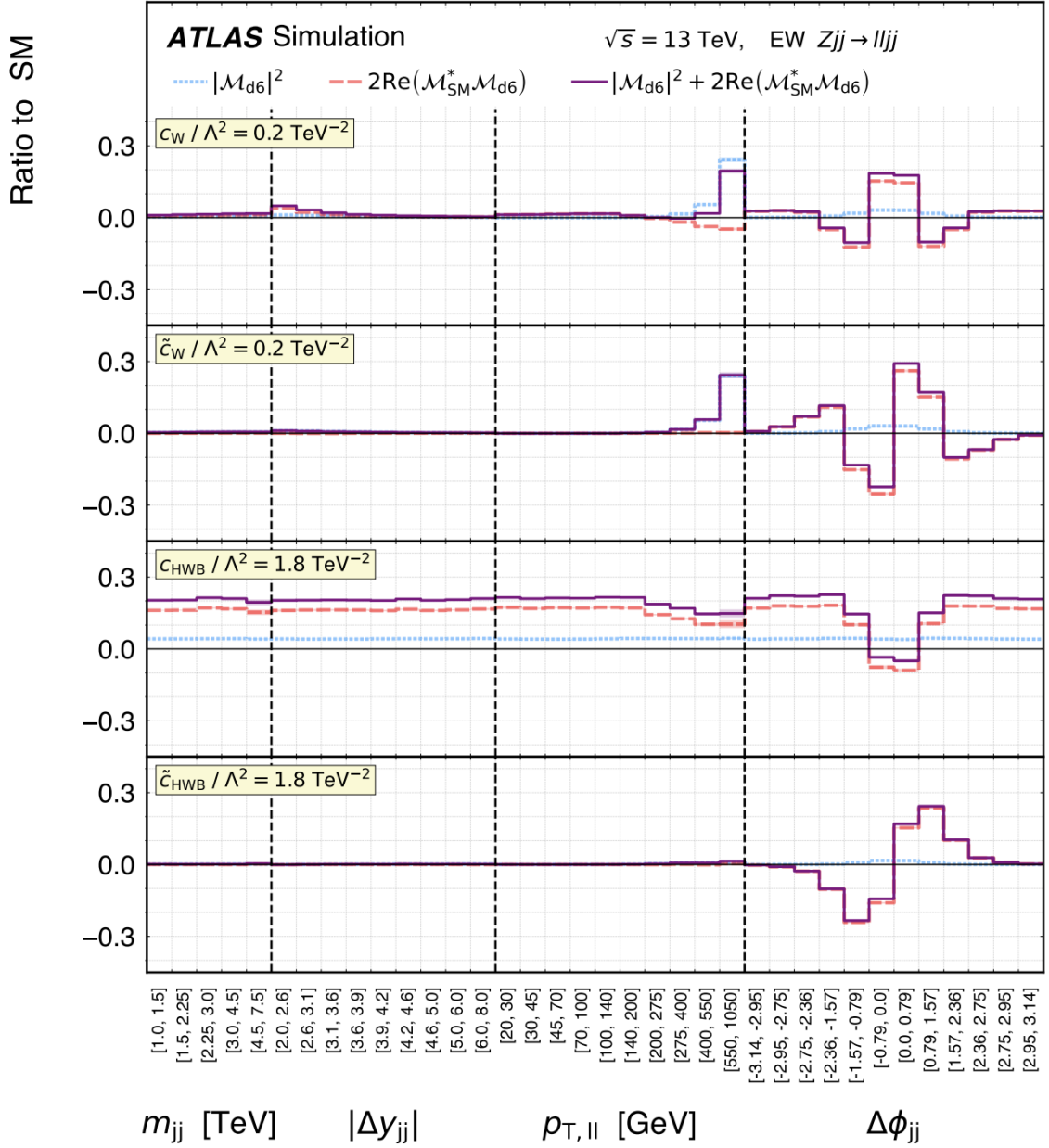


FIGURE 7.1: Impact of the $\mathcal{O}_W, \mathcal{O}_{\tilde{W}}, \mathcal{O}_{HWB}$ and $Q_{\tilde{H}\tilde{W}B}$ operators on the EW- Zjj differential cross sections. The expected contributions from the pure dimension-six term ($|\mathcal{M}_{d6}|^2$) and from the interference between the SM and dimension six amplitudes ($2\text{Re}(\mathcal{M}_{SM}^* \mathcal{M}_{d6})$) are shown relative to the pure SM prediction and represented as dotted and dashed lines, respectively. The total contribution to the EW- Zjj cross section is shown as a solid line.

that the ratio is not flat as a function of $\Delta\phi_{jj}^{\text{signed}}$, therefore a residual shape correction was derived. Figure 7.3 shows the quadratic fit to the normalized MADGRAPH5+PY8 / HERWIG+VBFNLO ratio, designed to parameterize the shape difference. Three eigenvariations are obtained by diagonalising the covariance matrix of the three fit parameters, which have uncertainties determined from the statistical variation of the MADGRAPH5+PY8 sample. The 68% and 95% confidence intervals shown in this figure are determined from pseudo-experiments (*toys*) performed by sampling the eigenvariations. The impact of these corrections on the expected limits is $\mathcal{O}(10\%)$ from the k factor and $\mathcal{O}(2\%)$ from the shape correction. A 9% theoretical uncertainty (described in Section 6.4.2) is assigned to the HERWIG+VBFNLO SM prediction, this is considered along with the statistical uncertainty of the prediction during EFT limit setting.

7.2 Statistical methods

This section will outline the statistical method employed to set limits on the Wilson coefficients associated to the four EFT operators using the EW- Zjj measurement as a function of $\Delta\phi_{jj}^{\text{signed}}$. All uncertainties are treated as Gaussian. Experimental systematic uncertainties and the strong- Zjj generator choice uncertainty are profiled as nuisance parameters. Theory uncertainties on the SM and EFT predictions are not profiled. The HERWIG+VBFNLO prediction is used for the SM contribution, since it was found to be in good agreement with the unfolded result. For a particular

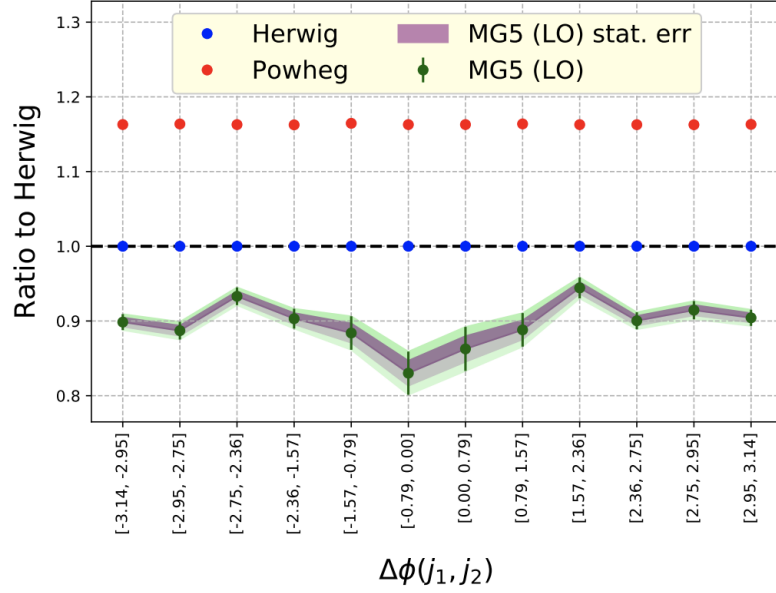


FIGURE 7.2: Comparison of the POWHEG+PY8, HERWIG+VBFNLO and MADGRAPH5+PY8 SM EW- Zjj predictions as a function of $\Delta\phi_{jj}^{\text{signed}}$. The ratio of each sample to the nominal HERWIG+VBFNLO prediction is presented.

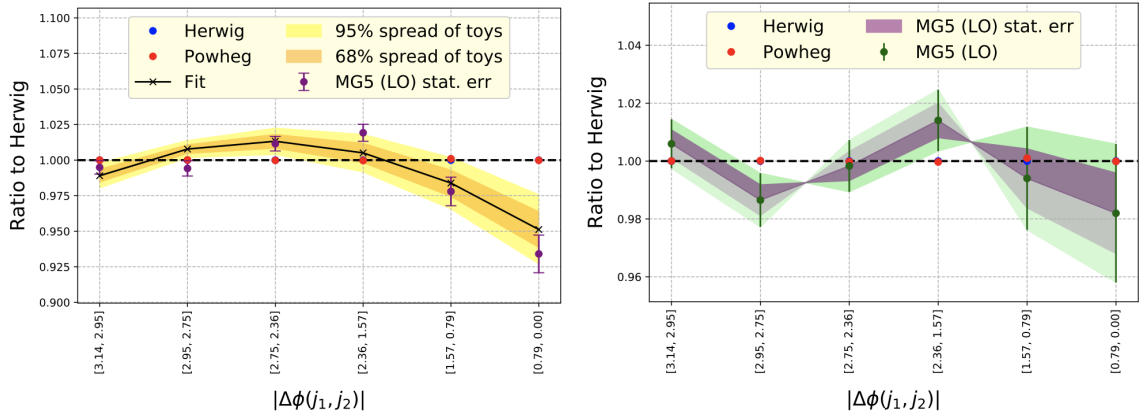


FIGURE 7.3: Left: Normalized MADGRAPH5+PY8 / HERWIG+VBFNLO ratio with a quadratic fit to the shape. The orange and yellow band are the uncertainty on the fit parameters, derived from the statistical uncertainty via pseudo-experiments (toys). Right: The same ratio after this shape correction is applied.

Wilson coefficient c_k the χ^2 function dependent on c_k and on the systematic NPs θ is

$$\chi^2(c_k, \theta) = \left(\mathbf{x} + \sum_i \theta_i \cdot \mathbf{e}_\theta - \mathbf{p}(c_k) \right)^T \text{Cov}^{-1} \left(\mathbf{x} + \sum_i \theta_i \cdot \mathbf{e}_\theta - \mathbf{p}(c_k) \right), \quad (7.4)$$

where $\mathbf{x}$ is the vector of measurements and $\mathbf{p}(c_k)$ is the prediction for a given value of c_k . Cov is the sum of the statistical covariances on $\mathbf{x}$ and $\mathbf{p}$, and θ are nuisance parameters where each component θ_i represents the fractional pull on a single experimental systematic uncertainty with an absolute magnitude of $\mathbf{e}_\theta$. The likelihood function is then

$$\mathcal{L}(c_k, \theta) = \frac{1}{\sqrt{(2\pi)^n \cdot |\text{Cov}|}} \cdot e^{-\frac{1}{2}\chi^2(c_k, \theta)} \cdot \prod_i \mathcal{G}(\theta_i), \quad (7.5)$$

where $\mathcal{G}$ represents a unit Gaussian distribution evaluated at θ_i and n is the length of the measurement vector $\mathbf{x}$.

The likelihood function is used to set the 95% confidence limit on the Wilson coefficient by construction a test statistic [115] $q(c_k)$ defined as

$$q(c_k) = -2 \log \frac{\mathcal{L}(c_k, \hat{\theta})}{\mathcal{L}(\hat{c}_k, \hat{\theta})} \quad (7.6)$$

where $\hat{c}_k$ and $\hat{\theta}$ are the maximum likelihood estimators and $\hat{\theta}$ are the values of θ that maximise the likelihood at a given value of c_k . The test statistic is distributed according to a χ^2 distribution with one degree of freedom following from Wilks' theorem [116]. The confidence limits are constructed by excluding values of c_k for

which

$$\int_0^{q(c_k)} \chi^2 (\text{d.o.f.} = 1) dq > 95 \%. \quad (7.7)$$

7.3 Effective field theory limits

Table 7.1 shows the expected and observed 95% confidence intervals on the Wilson coefficients associated to the four dimension six EFT operators. The confidence intervals were determined with and without considering the pure dimension six contribution to the EFT prediction. As expected from the observations of the spectra of Figure 7.1, the confidence intervals are only slightly affected by the inclusion or exclusion of this component. The greatest expected sensitivity is for c_W and $\tilde{c}_W$, the size of the confidence intervals in c_{HWB} and $\tilde{c}_{HWB}$ are roughly an order of magnitude larger. The p-values presented in Table 7.1 quantify the compatibility with the SM such that 100% denotes perfect agreement.

The measurement of the Wilson coefficient $\tilde{c}_{HWB}$, associated to the $\tilde{\mathcal{O}}_{HWB}$ operator, has a 95% confidence interval of [0.23, 2.34]. This interval does not include the SM (i.e. $\tilde{c}_{HWB} = 0$). The compatibility between the measurement and the SM is associated to a p -value of 1.6%. This confidence limit, and the others presented for each Wilson coefficient are local-significances. Since there are four different signal hypotheses under test using the same data there is a significant *look-elsewhere-effect*. Pseudo-experiments were once again employed to test the global significance of the result. From 10,000 pseudo-experiments it was observed that 6.2% exceeded the 1.6% significance threshold, therefore the global significance of the local 1.6%

Wilson coefficient	Includes $ \mathcal{M}_{d6} ^2$	95 % confidence interval [TeV ⁻²]		p -value (SM)
		Expected	Observed	
c_W	no	[−0.30, 0.30]	[−0.19, 0.41]	45.9 %
	yes	[−0.31, 0.29]	[−0.19, 0.41]	43.2 %
$\tilde{c}_W$	no	[−0.12, 0.12]	[−0.11, 0.14]	82.0 %
	yes	[−0.12, 0.12]	[−0.11, 0.14]	81.8 %
c_{HWB}	no	[−2.45, 2.45]	[−3.78, 1.13]	29.0 %
	yes	[−3.11, 2.10]	[−6.31, 1.01]	25.0 %
$\tilde{c}_{HWB}$	no	[−1.06, 1.06]	[0.23, 2.34]	1.7 %
	yes	[−1.06, 1.06]	[0.23, 2.35]	1.6 %

TABLE 7.1: Expected and observed 95% confidence intervals of the four Wilson coefficients, using fits to the $\Delta\phi_{jj}^{\text{signed}}$ observable. Results are presented with and without the quadratic contribution to the EFT prediction. The p -value quantifying the compatibility with the SM hypothesis is also shown for each limit.

$\tilde{\mathcal{O}}_{HWB}$ result is 6.2%. There is a reasonable chance that this high local significance is caused by the statistical fluctuations of the data that happen to be similar to the EFT MC prediction for the $\tilde{\mathcal{O}}_{HWB}$ operator.

An additional validation of the overall analysis procedure is presented in Appendix C. This is the so-called *injection test* where the EFT signal templates are used to inject an additional component into the EW- Zjj signal according to the boundary values of the 95% confidence interval of the Wilson coefficients. The purpose of this validation is to show that the EW- Zjj extraction procedure is sensitive to the expected EFT signal.

Chapter 8

Conclusion

The differential cross section of the electroweak production of a Z boson in association with two jets (EW- Zjj) has been measured for the first time. This measurement was performed using the pp collision data collected by the ATLAS experiment at a centre-of-mass energy of $\sqrt{s} = 13$ TeV and with an integrated luminosity of 139 fb^{-1} . This EW- Zjj process is sensitive to the vector-boson fusion process and therefore probes the WWZ interaction of the SM.

The cross section was measured in the $Z \rightarrow \ell^+ \ell^-$ decay channel as a function of four observables: the invariant mass of the two jet system, the rapidity interval spanned by the two jets, the signed azimuthal angle between the two jets, and the transverse momentum of the Z boson. The observed total fiducial cross section of EW- Zjj is 37.14 ± 3.47 (stat.) ± 5.79 (syst.) fb. An advanced methodology was developed to derive a data driven correction to correct the dominant strong- Zjj background in the signal region in order to extract the EW- Zjj signal. The observed results are corrected for detector inefficiency and resolution using an iterative Bayesian unfolding procedure and are compared to state-of-the-art predictions. The unfolded differential EW- Zjj cross section measurements are presented in

Figure 8.1, theoretical predictions from HERWIG+VBFNLO, POWHEG+PY8 and SHERPA are overlaid for comparison. The data results favour the HERWIG+VBFNLO prediction, POWHEG+PY8 over predicts the data at high dijet invariant mass, for large dijet rapidity intervals, and at intermediate values of dilepton transverse momentum. Sherpa predicts too small a cross-section across the measured phase space. Differential cross-section measurements for inclusive Zjj production are also provided in the signal and control regions used to extract the electroweak component. These strong- Zjj dominated measurements are in good agreement with the SM predictions from MG5_NLO+PY8 and SHERPA.

The differential cross sections were used to test deviations from the SM attributed to dimension-six corrections in the WWZ vertex by exploiting the sensitivity of the signed azimuthal angle between the two jets with an effective field theory approach. Overall, the results are consistent with the SM. The dimension-six operators are constrained primarily by the contribution to the cross section from the interference between the SM and dimension-six scattering amplitudes. The effective field theory limits placed on the anomalous weak boson self interactions constitute an important test of charge conjugation and parity in the SM.

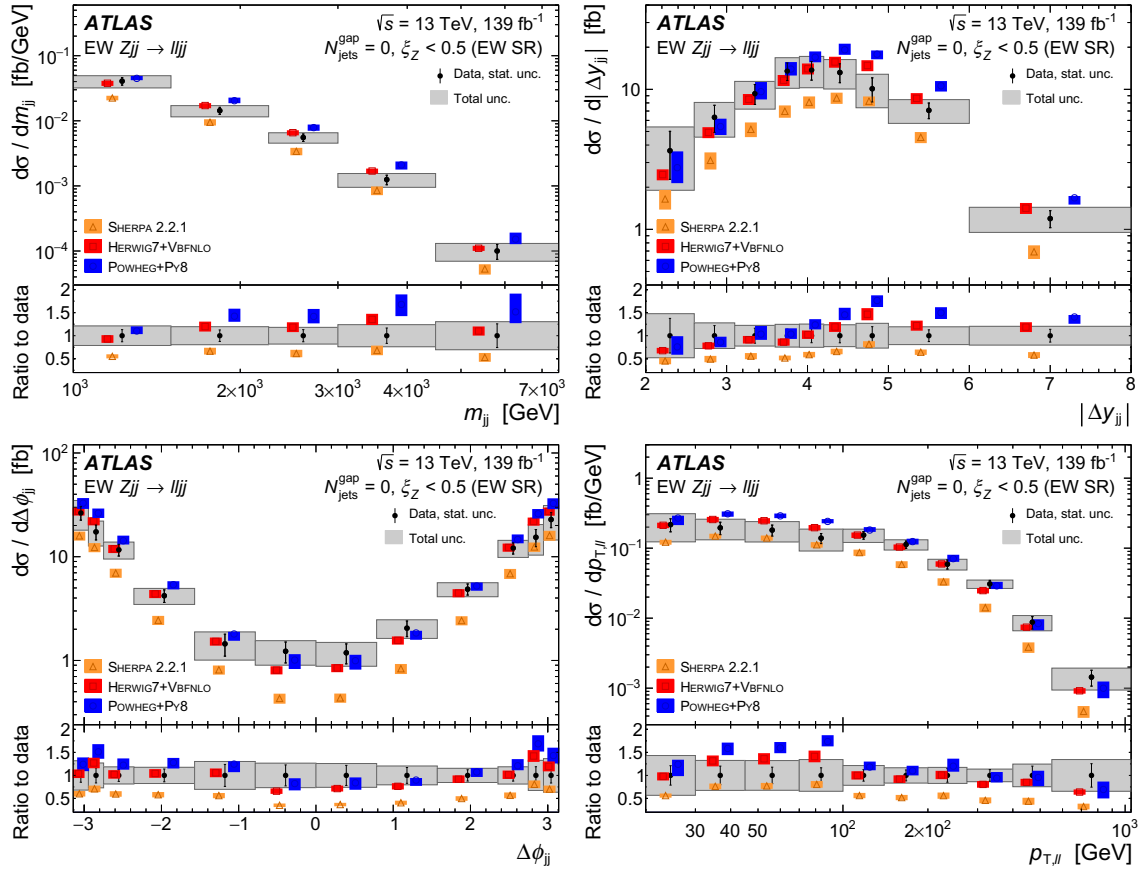


FIGURE 8.1: Differential cross sections for EW-Zjj production as a function of the four measurement variables. The unfolded data are shown as black points, with the statistical uncertainty represented by an error bar and the total uncertainty represented as a grey band. The data are compared with theoretical predictions produced by HERWIG + VBFNLO (red points), POWHEG+PY8 (blue points) and SHERPA (orange points). Uncertainty bands are shown for the three theoretical predictions. Each theory prediction is slightly offset from the bin centre to avoid overlap.

Appendix A

Particle-level simulation reweighting

As described in Section 4.1.2, two samples, HERWIG+VBFNLO EW- Zjj and MG5_NLO+PY8 strong- Zjj , exist only at particle level. Approximate detector-level predictions for these samples are obtained by tuning the MC event weights (reweighting) of another MC sample (one that does have a detector-level prediction) to match kinematic distributions of the particle level only prediction. Samples where such a particle-level reweighting has been applied are labelled with a prime symbol. The detector-level prediction of HERWIG+VBFNLO EW- Zjj is obtained by reweighting POWHEG+PY8 EW- Zjj to match HERWIG+VBFNLO at particle level, this sample is referred to as HERWIG+VBFNLO'. Similarly, MG5_NLO+PY8' is obtained by reweighting MG5+PY8 strong- Zjj at particle level.

The simulated (Asimov) datasets used in Section 6.3 were created via a similar procedure, applying a reweighting function to the strong- Zjj MC at particle level such that the prediction matches the data at detector level in the control regions.

The modelling of the EW- Zjj process is generally consistent between the HERWIG+VBFNLO and POWHEG+PY8 generators in term of the shape of the distributions of the four measurement variables. The difference between the samples

is essentially a difference in overall normalization, the event yield from HERWIG+VBFNLO is consistently 0.85 that of POWHEG+PY8 in nearly all bins of the four measurement variables. This is shown for the m_{jj} distribution in Figure A.1 in the VBF topology region, the ratio of HERWIG+VBFNLO and SHERPA to POWHEG+PY8 EW- Zjj is also shown in each of control regions and in the signal region. The ratio shown in the VBF topology region (the sum of CRa, CRb, CRc and the SR) is very similar to that of the SR since it is where the vast majority of EW- Zjj events reside. The SHERPA sample is included for comparison, the normalization of this sample is roughly half that of the other two samples as a result of a non-optimal setting of the colour flow [74].

The reweighting of POWHEG+PY8 to obtain HERWIG+VBFNLO' is performed by applying an event by event weight to the POWHEG+PY8 events equal to the ratio of the two samples in the particle-level m_{jj} bin in the fiducial control or signal region to which the event belongs. These values are shown in the ratios of Figure A.1.

The modelling of strong- Zjj varies more between generators than that of EW- Zjj . As such the particle-level reweighting of MG5+PY8 to obtain MG5_NLO+PY8' is more complex. A reweighting function was determined by iteration, following a procedure of checking the modelling in a variety of variables important to the analysis. The final reweighting function is parametrized in terms of the following observables, all at particle level: m_{jj} , $|\Delta y_{jj}|$, $p_T^{\ell\ell jj}$, p_T^{j2} , ξ_Z and $N_{\text{jets}}^{\text{gap}}$. The iteration involves starting with a reweighting function parametrized in terms of a single observable, this will be insufficient and leave a residual difference in other observables. This agreement is tested at detector level, and observables are added to the reweighting function until the residual correction in a variety of variables

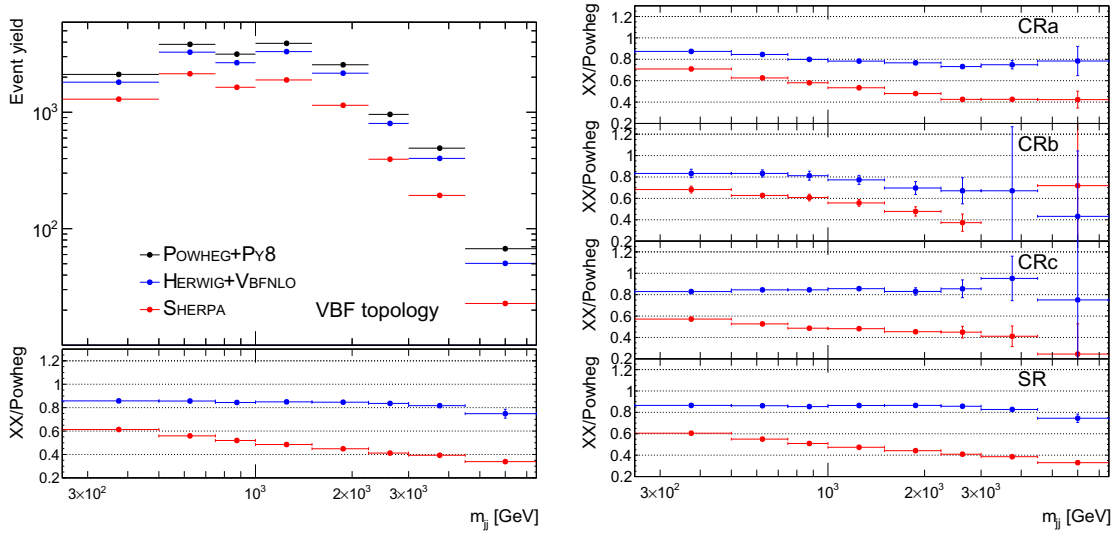


FIGURE A.1: Comparison of EW- Zjj samples at particle level as a function of m_{jj} . The event yield of the three EW- Zjj sample is shown on the left with the ratio to POWHEG+PY8 of the other two samples. The same ratios are shown on the right, this time in each of the control regions and in the signal region.

has been reduced. The final reweighted sample MG5_NLO+PY8' is shown in Figure A.2, the agreement to the data is greatly improved in the reweighted sample. The particle-level reweighting used to obtain the MG5_NLO+PY8' sample is implemented with the following C++ function.

```
double getFxFxWeight(double truth_pT_Zjj, double truth_pT_j2, double truth_Dy_jj,
double truth_m_jj, double truth_Z_centrality, int truth_NgapJets){
return (truth_pT_Zjj>0 ? (truth_pT_Zjj>75 ? -1.7116+4.1232*pow(truth_pT_Zjj,-0.10164)
: 0.34886+0.021449*truth_pT_Zjj-0.0001798*pow(truth_pT_Zjj,2)) : 1.0)

*(truth_pT_j2>80 ? 7.3226*pow(truth_pT_j2,-0.417) : 1.0)

*(truth_Dy_jj>2 ? 0.64661+0.439815*truth_Dy_jj-0.14944*pow(truth_Dy_jj,2)+0.0137815*pow(truth_Dy_jj,3) : 1.0)

*(truth_m_jj>250 ? 1.01055-2.78831e-05*truth_m_jj : 1.0)

*(truth_m_jj>1000 ? ( truth_pT_Zjj>140 ? 1.179321 : 0.743259+0.0062306*truth_pT_Zjj-2.22563e-05*pow(truth_pT_Zjj,2) ) : 1.0)

*(truth_Z_centrality>0 ? 0.921+0.3483*pow(truth_Z_centrality,2) : 1.0)

*(truth_NgapJets==0 ? 0.95 : 1.07);
}
```

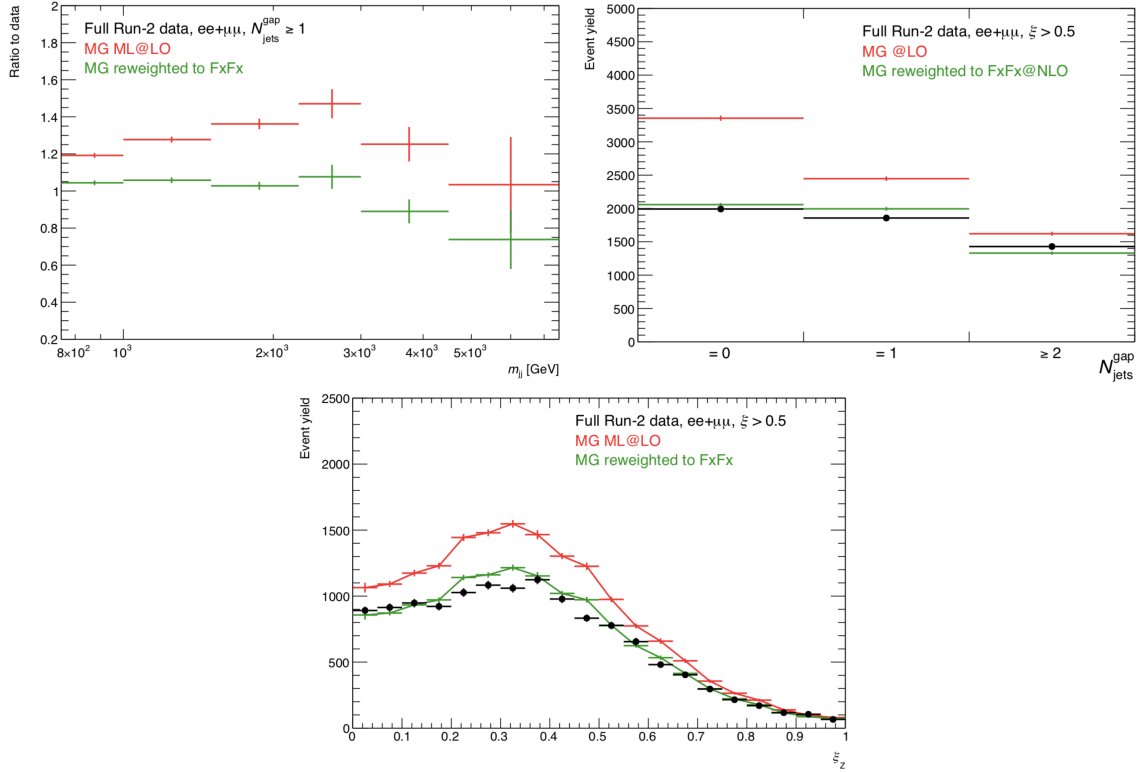


FIGURE A.2: Result of the reweighting of MG5+PY8 to obtain MG5_NLO+PY8' strong- Zjj . The ratio to the data is shown before (red) and after (green) reweighting as a function of m_{jj} (top left). MG5_NLO+PY8 is labelled as "FxFx" in this figure, this refers to the matching and merging scheme used in the NLO simulation. The distribution of $N_{\text{jets}}^{\text{gap}}$ (top right) and ξ_Z are also shown before and after reweighting and the agreement to the Run II data is greatly improved by the reweighting.

Appendix B

Additional unfolded results material

This appendix includes additional differential cross section and uncertainty figures for the inclusive measurements performed in the three control regions, presented in Figures B.1–B.6. The unfolded differential cross sections for both EW- Zjj production and inclusive Zjj production are documented in tabular form in Tables B.1–B.4.

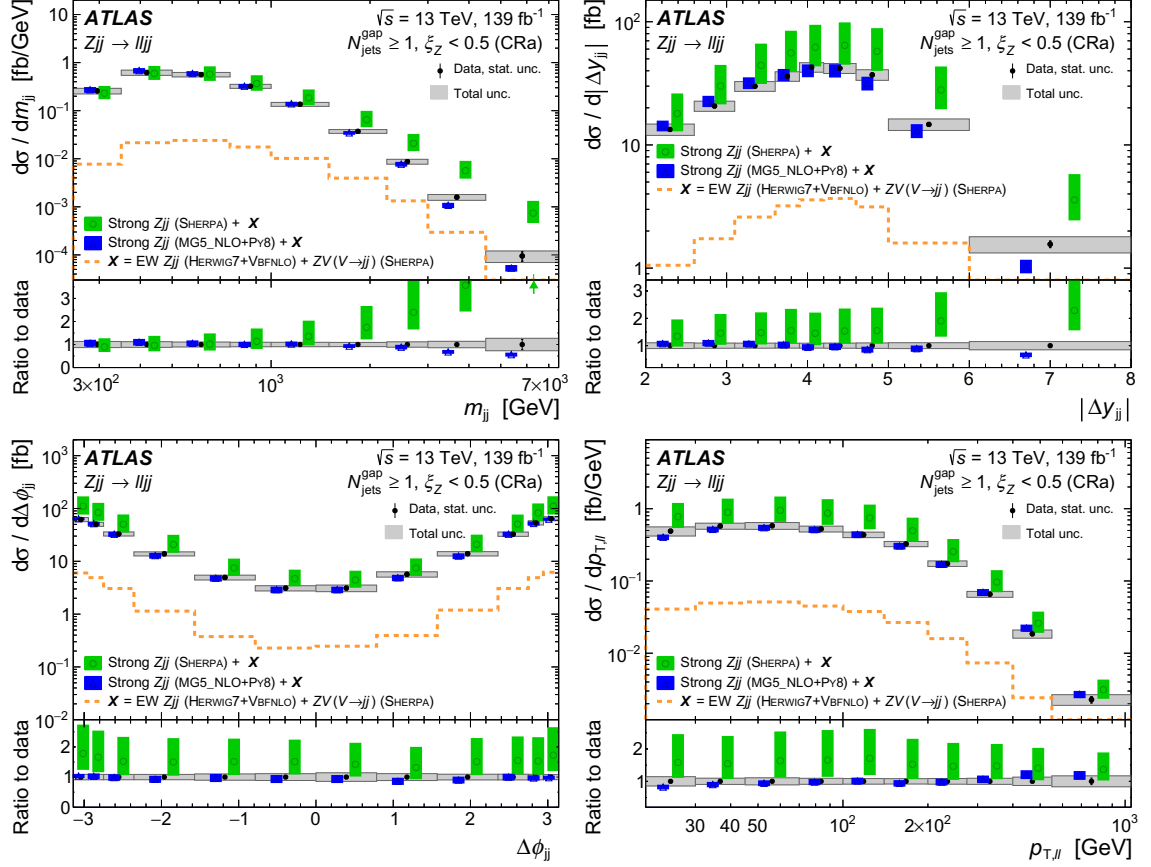


FIGURE B.1: Differential cross-sections measured in CRa for inclusive Zjj production as a function of the four measurement variables in the signal region. The unfolded data are shown as black points, with the statistical uncertainty represented by an error bar and the total uncertainty represented as a grey band. The data are compared with theoretical predictions constructed from different strong- Zjj predictions provided by SHERPA (green) and MG5_NLO+PY8 (blue). Uncertainty bands are shown for the two theoretical predictions. Each theory prediction is slightly offset from the bin centre to avoid overlap.

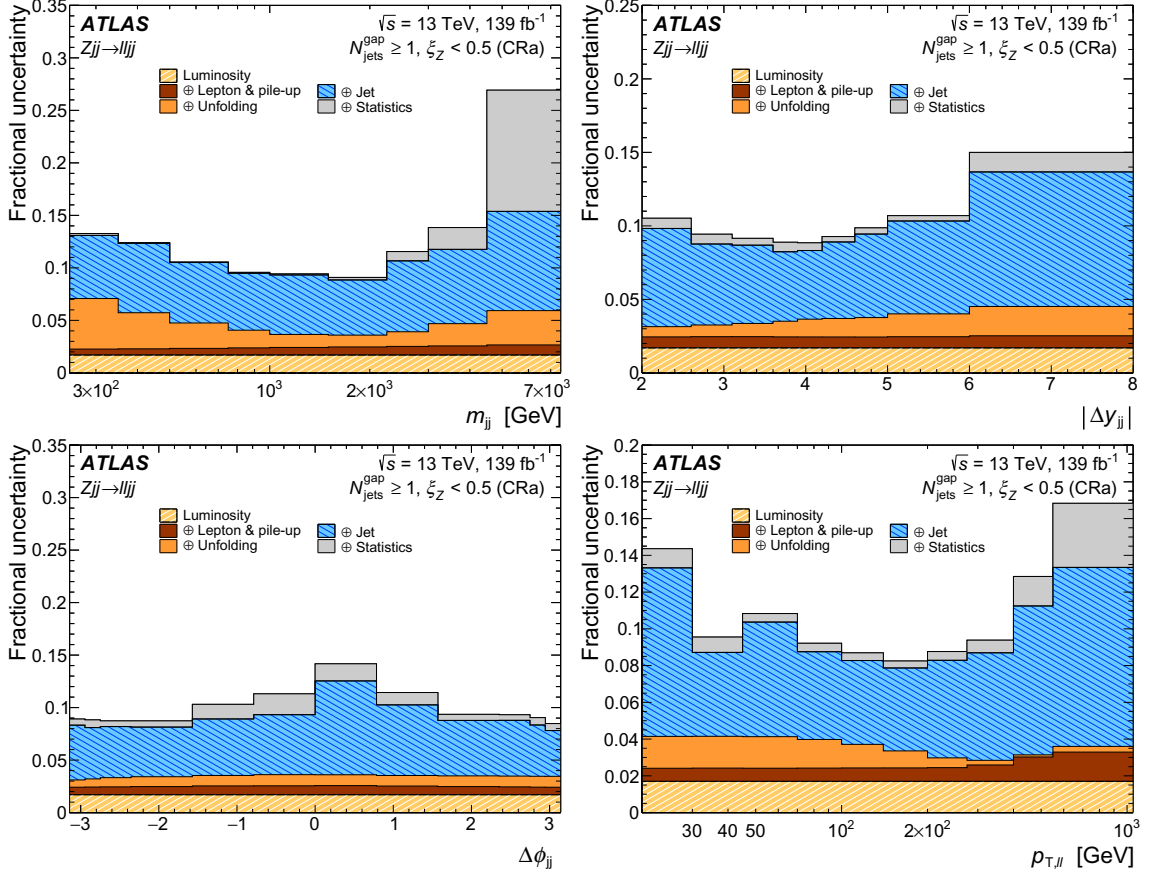


FIGURE B.2: Fractional uncertainty in the inclusive Zjj measurement as a function of the four measurement variables in CRa. Uncertainty sources are grouped in categories that are added in quadrature (denoted $\oplus$) to give the total uncertainty.

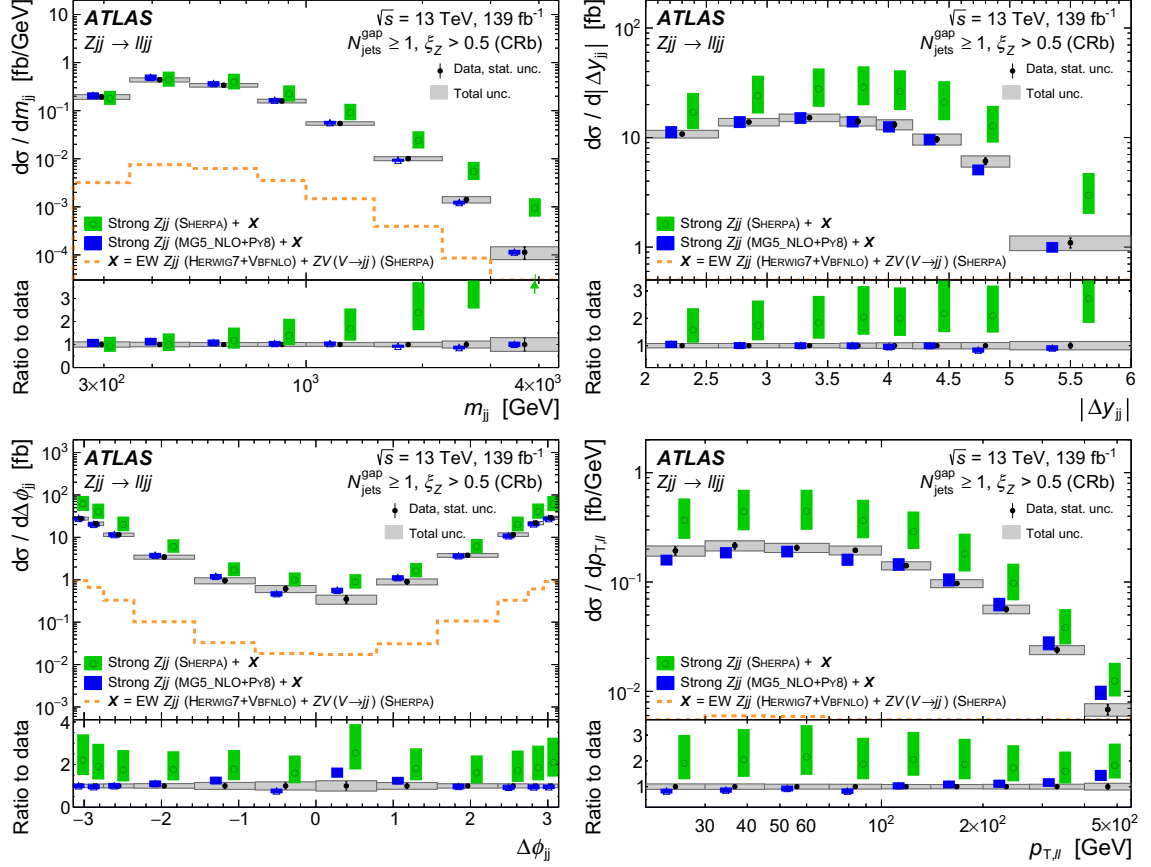


FIGURE B.3: Differential cross-sections measured in CRb for inclusive Zjj production as a function of the four measurement variables in the signal region. The unfolded data are shown as black points, with the statistical uncertainty represented by an error bar and the total uncertainty represented as a grey band. The data are compared with theoretical predictions constructed from different strong- Zjj predictions provided by SHERPA (green) and MG5_NLO+PY8 (blue). Uncertainty bands are shown for the two theoretical predictions. Each theory prediction is slightly offset from the bin centre to avoid overlap.

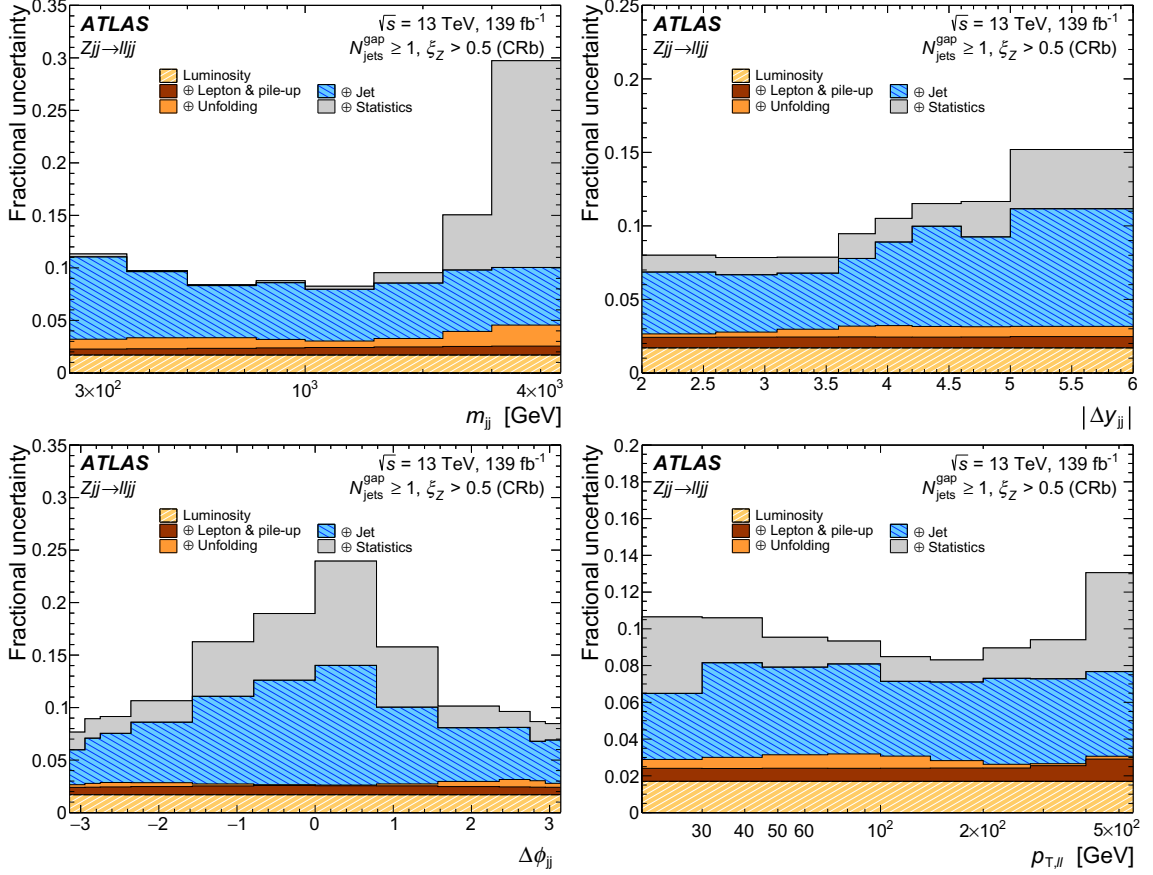


FIGURE B.4: Fractional uncertainty in the inclusive Zjj measurement as a function of the four measurement variables in CRb. Uncertainty sources are grouped in categories that are added in quadrature (denoted $\oplus$) to give the total uncertainty.

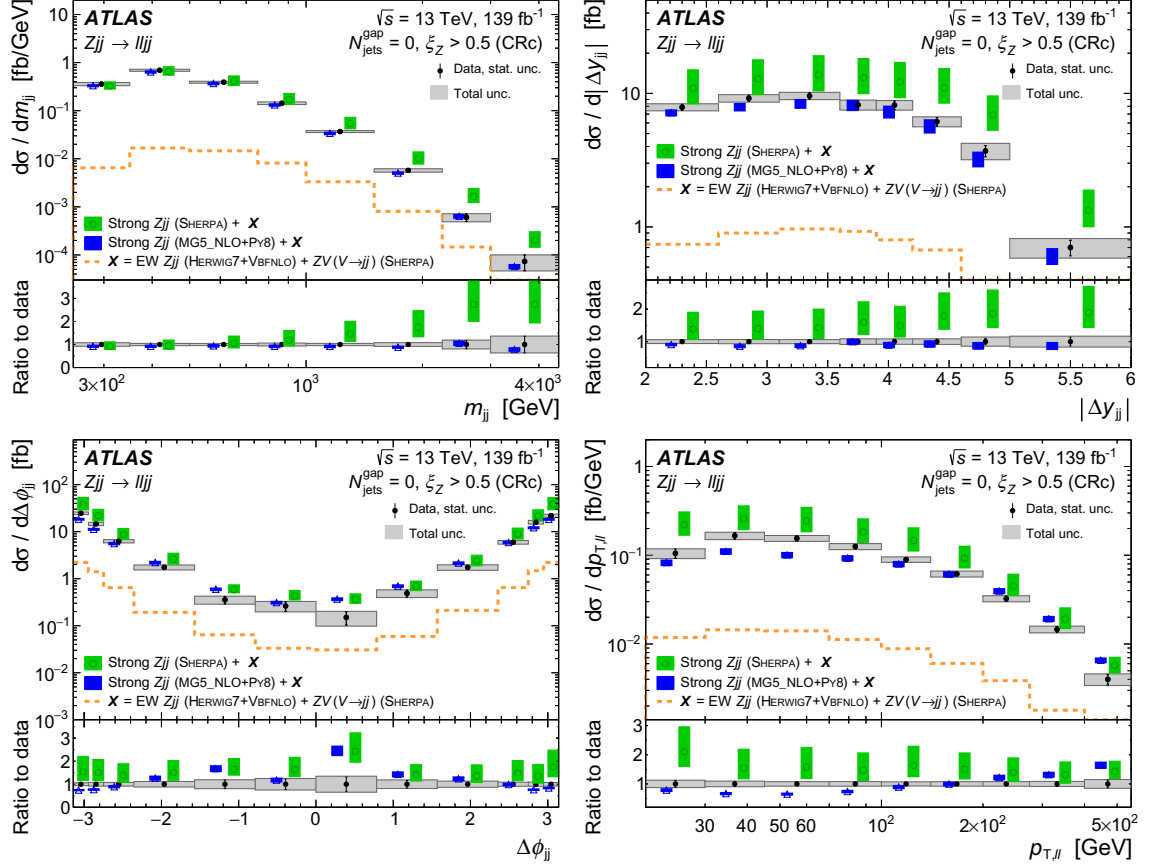


FIGURE B.5: Differential cross-sections measured in CRc for inclusive Zjj production as a function of the four measurement variables in the signal region. The unfolded data are shown as black points, with the statistical uncertainty represented by an error bar and the total uncertainty represented as a grey band. The data are compared with theoretical predictions constructed from different strong- Zjj predictions provided by SHERPA (green) and MG5_NLO+PY8 (blue). Uncertainty bands are shown for the two theoretical predictions. Each theory prediction is slightly offset from the bin centre to avoid overlap.

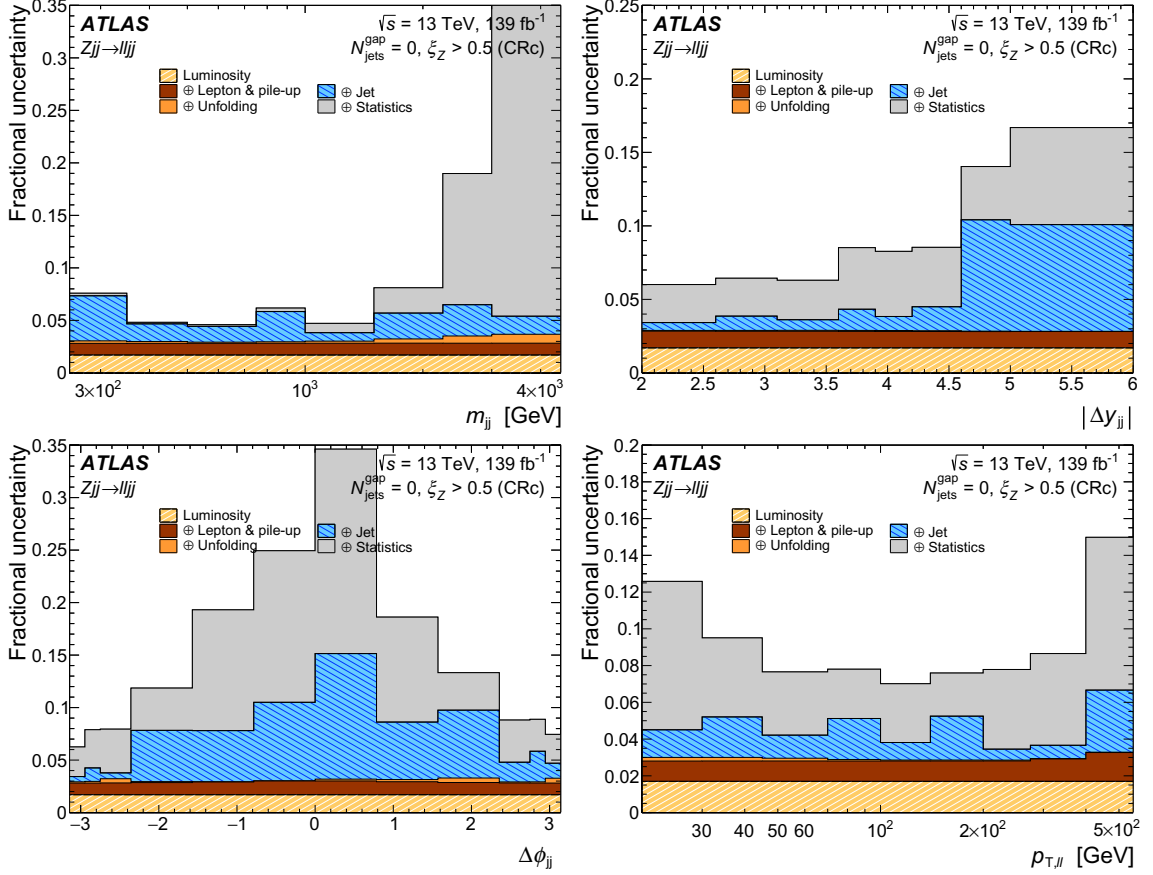


FIGURE B.6: Fractional uncertainty in the inclusive Zjj measurement as a function of the four measurement variables in CRc. Uncertainty sources are grouped in categories that are added in quadrature (denoted $\oplus$) to give the total uncertainty.

EW Zjj SR, m_{jj} cross-section measurements									
$d\sigma / dm_{jj}$ [ab/GeV]	-	-	-	-	41	14	5.5	1.3	0.10
Stat. unc. [%]	-	-	-	-	13	13	13	17	26
Gen. choice [%]	-	-	-	-	11	11	9.4	14	7.6
Theory syst. [%]	-	-	-	-	8.1	6.6	4.3	3.1	1.2
Jet syst. [%]	-	-	-	-	8.4	6.9	6.3	9.4	14
Unfolding syst. [%]	-	-	-	-	2.3	1.1	0.7	0.6	0.6
Other syst. [%]	-	-	-	-	2.0	2.0	2.3	2.2	3.0
Inclusive Zjj SR, m_{jj} cross-section measurements									
$d\sigma / dm_{jj}$ [ab/GeV]	510	1040	700	320	120	31	8.8	1.7	0.12
Stat. unc. [%]	1.6	1.0	0.9	1.3	1.5	2.3	4.5	7.2	21
Jet syst. [%]	5.2	3.8	3.3	3.6	3.6	3.5	4.1	6.6	15
Unfolding syst. [%]	2.3	1.6	0.9	0.6	0.5	0.4	0.5	0.6	0.6
Other syst. [%]	2.8	2.8	2.8	2.8	2.8	2.8	2.9	2.9	3.4
Inclusive Zjj CRa, m_{jj} cross-section measurements									
$d\sigma / dm_{jj}$ [ab/GeV]	250	610	560	320	130	37	8.7	1.6	0.10
Stat. unc. [%]	2.2	1.2	1.0	1.3	1.3	2.1	4.4	7.3	22
Jet syst. [%]	11	11	9.4	8.6	8.6	8.1	9.9	11	14
Unfolding syst. [%]	6.7	5.3	4.1	3.3	2.7	2.6	3.0	3.9	5.3
Other syst. [%]	2.3	2.3	2.3	2.4	2.4	2.5	2.5	2.6	2.8
Inclusive Zjj CRb, m_{jj} cross-section measurements									
$d\sigma / dm_{jj}$ [ab/GeV]	190	430	330	150	54	10	1.4	0.11	-
Stat. unc. [%]	2.5	1.4	1.2	1.8	2.2	4.2	11	28	-
Jet syst. [%]	11	9.0	7.6	8.0	7.4	7.9	9.0	8.9	-
Unfolding syst. [%]	2.3	2.4	2.4	2.1	1.8	2.1	3.0	3.8	-
Other syst. [%]	2.3	2.3	2.3	2.4	2.4	2.5	2.6	2.6	-
Inclusive Zjj CRc, m_{jj} cross-section measurements									
$d\sigma / dm_{jj}$ [ab/GeV]	350	690	390	140	37	5.7	0.60	0.07	-
Stat. unc. [%]	1.9	1.2	1.2	2.0	2.7	5.8	18	36	-
Jet syst. [%]	6.7	3.6	3.3	5.0	2.3	4.7	5.5	4.0	-
Unfolding syst. [%]	1.2	1.0	0.8	0.9	1.1	1.6	2.1	2.3	-
Other syst. [%]	2.8	2.8	2.8	2.8	2.8	2.9	2.9	3.1	-
Low bin edge [TeV]	0.25	0.35	0.50	0.75	1.0	1.5	2.2	3.0	4.5
High bin edge [TeV]	0.35	0.50	0.75	1.0	1.5	2.2	3.0	4.5	7.5

TABLE B.1: Differential cross-section measurements for EW- Zjj production and inclusive Zjj production as a function of m_{jj} . The EW- Zjj measurements are performed in the signal region. The inclusive Zjj measurements are performed in the signal region and three control regions. The different sources of uncertainty are grouped together for each measurement, with a more granular breakdown of each uncertainty available in HEPDATA.

EW Zjj SR, Δy_{jj} cross-section measurements									
$d\sigma / d \Delta y_{jj} $ [fb]	3.6	6.3	9.3	13	14	13	10	7.1	1.2
Stat. unc. [%]	38	22	16	14	15	15	20	13	14
Gen. choice [%]	7.5	4.0	9.8	17	17	14	10	12	10
Theory syst. [%]	18	11	8.5	6.7	8.0	7.5	7.5	3.6	1.8
Jet syst. [%]	21	10.0	6.9	7.0	6.4	8.3	13	6.2	11
Unfolding syst. [%]	8.4	6.3	3.9	2.8	2.3	1.5	1.1	0.7	0.1
Other syst. [%]	2.6	2.3	2.2	2.4	2.1	2.0	2.0	1.9	2.2
Inclusive Zjj SR, Δy_{jj} cross-section measurements									
$d\sigma / d \Delta y_{jj} $ [fb]	13	18	25	32	36	36	31	14	1.6
Stat. unc. [%]	3.9	3.8	3.2	3.5	3.3	2.9	3.2	3.0	6.3
Jet syst. [%]	2.7	2.1	2.2	2.6	2.9	5.7	4.5	4.2	7.7
Unfolding syst. [%]	0.7	0.6	0.5	0.4	0.3	0.3	0.4	0.4	0.3
Other syst. [%]	2.9	2.9	2.9	2.9	2.9	2.8	2.9	2.8	2.9
Inclusive Zjj CRa, Δy_{jj} cross-section measurements									
$d\sigma / d \Delta y_{jj} $ [fb]	13	21	30	36	43	42	37	15	1.6
Stat. unc. [%]	3.8	3.5	2.9	3.3	3.0	2.6	2.9	2.8	6.2
Jet syst. [%]	9.3	8.1	8.0	7.5	7.5	8.1	8.7	9.5	13
Unfolding syst. [%]	2.0	2.1	2.3	2.5	2.7	2.8	2.9	3.2	3.7
Other syst. [%]	2.5	2.5	2.5	2.5	2.5	2.5	2.5	2.5	2.6
Inclusive Zjj CRb, Δy_{jj} cross-section measurements									
$d\sigma / d \Delta y_{jj} $ [fb]	11	14	15	14	13	9.6	6.1	1.1	-
Stat. unc. [%]	4.1	4.1	4.0	5.4	5.6	5.8	7.1	10	-
Jet syst. [%]	6.3	6.1	6.1	7.1	8.3	9.5	8.7	11	-
Unfolding syst. [%]	1.1	1.3	1.7	2.0	2.1	2.0	2.0	2.0	-
Other syst. [%]	2.4	2.4	2.5	2.5	2.5	2.5	2.5	2.6	-
Inclusive Zjj CRc, Δy_{jj} cross-section measurements									
$d\sigma / d \Delta y_{jj} $ [fb]	7.9	9.2	9.6	8.2	8.2	6.1	3.7	0.70	-
Stat. unc. [%]	4.9	5.2	5.2	7.3	7.3	7.3	9.4	13	-
Jet syst. [%]	1.8	2.6	2.2	3.2	2.5	3.5	10	9.7	-
Unfolding syst. [%]	0.5	0.6	0.6	0.6	0.6	0.5	0.4	0.2	-
Other syst. [%]	2.9	2.9	2.9	2.9	2.9	2.9	3.1	3.4	-
Low bin edge	2.0	2.6	3.1	3.6	3.9	4.2	4.6	5.0	6.0
High bin edge	2.6	3.1	3.6	3.9	4.2	4.6	5.0	6.0	8.0

TABLE B.2: Differential cross-section measurements for EW- Zjj production and inclusive Zjj production as a function of $|\Delta y_{jj}|$. The EW- Zjj measurements are performed in the signal region. The inclusive Zjj measurements are performed in the signal region and three control regions. The different sources of uncertainty are grouped together for each measurement, with a more granular breakdown of each uncertainty available in HEPDATA.

EW Zjj SR, $\Delta\phi_{jj}$ cross-section measurements												
$d\sigma / d\Delta\phi_{jj}$ [fb]	26	17	12	4.2	1.4	1.2	1.2	2.0	4.9	12	15	23
Stat. unc. [%]	15	16	13	14	24	23	22	17	12	12	19	17
Gen. choice [%]	25	17	7.4	4.5	-14	-10	-11	-6.9	2.5	8.8	22	30
Theory syst. [%]	7.8	9.5	8.2	5.1	4.6	4.7	4.3	3.2	4.3	8.2	12	8.2
Jet syst. [%]	8.9	8.0	6.6	7.9	11	7.0	7.6	7.5	7.4	6.3	8.7	9.9
Unfolding syst. [%]	2.6	2.4	2.2	2.1	1.8	1.1	0.6	0.8	1.7	2.8	3.4	3.9
Other syst. [%]	2.6	2.1	2.1	2.1	2.3	2.6	2.5	2.3	2.1	2.1	2.2	2.7
Inclusive Zjj SR, $\Delta\phi_{jj}$ cross-section measurements												
$d\sigma / d\Delta\phi_{jj}$ [fb]	71	49	29	10	3.3	2.1	2.0	3.8	11	29	48	69
Stat. unc. [%]	3.1	3.6	3.3	3.8	6.4	8.1	8.6	6.1	3.7	3.3	3.8	3.3
Jet syst. [%]	2.0	5.1	3.6	4.2	4.4	8.2	6.6	8.0	4.6	3.4	4.0	2.6
Unfolding syst. [%]	0.4	0.4	0.3	0.3	0.3	0.1	0.1	0.1	0.2	0.4	0.5	0.6
Other syst. [%]	2.8	2.8	2.9	2.9	3.0	3.1	3.1	3.0	2.9	2.9	2.9	2.8
Inclusive Zjj CRa, $\Delta\phi_{jj}$ cross-section measurements												
$d\sigma / d\Delta\phi_{jj}$ [fb]	62	50	33	14	4.9	3.1	3.1	5.7	14	33	54	64
Stat. unc. [%]	3.2	3.5	3.0	3.2	5.2	6.4	6.6	5.1	3.2	3.1	3.5	3.3
Jet syst. [%]	7.7	7.4	7.5	7.4	8.2	8.6	12	9.6	8.1	8.1	7.6	7.0
Unfolding syst. [%]	2.0	2.1	2.3	2.4	2.5	2.6	2.5	2.5	2.5	2.5	2.5	2.5
Other syst. [%]	2.4	2.5	2.5	2.5	2.6	2.6	2.6	2.6	2.5	2.5	2.4	2.4
Inclusive Zjj CRb, $\Delta\phi_{jj}$ cross-section measurements												
$d\sigma / d\Delta\phi_{jj}$ [fb]	28	21	12	3.4	0.96	0.62	0.35	0.90	3.8	12	22	29
Stat. unc. [%]	4.8	5.4	5.2	6.3	12	14	19	12	6.1	5.2	5.4	4.9
Jet syst. [%]	5.4	6.5	7.0	8.1	11	12	14	9.7	7.5	7.5	6.1	6.3
Unfolding syst. [%]	1.1	1.4	1.5	1.4	1.0	0.6	0.6	1.0	1.6	2.0	1.8	1.4
Other syst. [%]	2.4	2.4	2.5	2.5	2.7	2.9	2.8	2.7	2.5	2.5	2.4	2.4
Inclusive Zjj CRc, $\Delta\phi_{jj}$ cross-section measurements												
$d\sigma / d\Delta\phi_{jj}$ [fb]	25	15	6.3	1.7	0.36	0.26	0.15	0.49	1.7	6.0	16	22
Stat. unc. [%]	5.2	6.6	7.0	8.9	18	23	31	16	9.1	7.4	6.7	5.8
Jet syst. [%]	1.8	3.0	2.0	7.3	7.2	10	15	8.0	9.2	3.8	5.1	3.4
Unfolding syst. [%]	0.8	0.8	0.7	0.4	0.2	0.2	0.2	0.2	0.3	0.6	0.7	0.7
Other syst. [%]	2.8	2.9	2.9	3.0	3.1	3.9	3.7	3.3	3.0	2.9	2.9	2.9
Low bin edge [$\pi/16$]	-16	-15	-14	-12	-8.0	-4.0	0.00	4.0	8.0	12	14	15
High bin edge [$\pi/16$]	-15	-14	-12	-8.0	-4.0	0.00	4.0	8.0	12	14	15	16

TABLE B.3: Differential cross-section measurements for EW- Zjj production and inclusive Zjj production as a function of $\Delta\phi_{jj}^{\text{signed}}$. The EW- Zjj measurements are performed in the signal region. The inclusive Zjj measurements are performed in the signal region and three control regions. The different sources of uncertainty are grouped together for each measurement, with a more granular breakdown of each uncertainty available in HEPDATA.

EW Zjj SR, $p_{T,\ell\ell}$ cross-section measurements										
$d\sigma / dp_{T,\ell\ell}$ [ab/GeV]	210	190	180	130	150	110	59	31	8.8	1.4
Stat. unc. [%]	21	20	18	17	12	12	15	13	20	25
Gen. choice [%]	36	19	22	24	14	6.9	4.2	-0.9	-12	-21
Theory syst. [%]	4.7	11	9.8	12	5.8	6.0	5.0	3.5	5.0	4.1
Jet syst. [%]	10	13	11	13	7.5	6.4	5.9	2.9	3.1	8.5
Unfolding syst. [%]	1.2	2.4	3.3	3.5	3.0	2.0	1.3	1.3	1.7	2.0
Other syst. [%]	3.5	2.1	2.1	2.2	2.1	2.0	2.1	2.2	3.0	3.4
Inclusive Zjj SR, $p_{T,\ell\ell}$ cross-section measurements										
$d\sigma / dp_{T,\ell\ell}$ [ab/GeV]	530	580	530	440	380	270	140	58	16	2.1
Stat. unc. [%]	5.2	4.0	3.3	3.2	2.9	2.8	3.2	3.9	7.0	11
Jet syst. [%]	5.3	6.9	4.1	4.2	3.6	3.6	2.2	1.4	2.5	5.6
Unfolding syst. [%]	0.4	0.4	0.5	0.5	0.4	0.3	0.2	0.2	0.3	0.3
Other syst. [%]	2.9	2.9	2.8	2.8	2.8	2.8	2.8	3.0	3.4	3.6
Inclusive Zjj CRa, $p_{T,\ell\ell}$ cross-section measurements										
$d\sigma / dp_{T,\ell\ell}$ [ab/GeV]	480	570	580	520	430	320	170	66	19	2.3
Stat. unc. [%]	5.4	3.9	3.1	2.9	2.7	2.5	2.9	3.5	6.2	10
Jet syst. [%]	13	7.7	9.5	7.8	7.4	7.1	7.7	8.2	11	13
Unfolding syst. [%]	3.4	3.4	3.4	3.2	2.8	2.3	1.6	1.1	0.8	0.7
Other syst. [%]	2.5	2.5	2.4	2.4	2.4	2.4	2.5	2.6	3.0	3.3
Inclusive Zjj CRb, $p_{T,\ell\ell}$ cross-section measurements										
$d\sigma / dp_{T,\ell\ell}$ [ab/GeV]	190	210	200	190	140	97	56	24	6.8	-
Stat. unc. [%]	8.4	6.8	5.3	4.7	4.6	4.3	5.2	5.9	11	-
Jet syst. [%]	5.8	7.6	7.3	7.4	6.4	6.5	6.8	6.8	7.0	-
Unfolding syst. [%]	1.6	1.8	2.0	2.1	1.9	1.5	1.0	0.7	0.8	-
Other syst. [%]	2.5	2.5	2.4	2.4	2.4	2.4	2.4	2.6	3.0	-
Inclusive Zjj CRc, $p_{T,\ell\ell}$ cross-section measurements										
$d\sigma / dp_{T,\ell\ell}$ [ab/GeV]	100	160	150	120	89	62	32	15	4.0	-
Stat. unc. [%]	12	8.0	6.4	5.9	5.9	5.5	7.0	7.8	13	-
Jet syst. [%]	3.4	4.3	3.0	4.2	2.5	4.4	1.9	2.2	5.8	-
Unfolding syst. [%]	1.1	1.1	0.9	0.7	0.5	0.5	0.5	0.3	0.3	-
Other syst. [%]	3.0	2.9	2.9	2.9	2.9	2.9	2.9	3.0	3.4	-
Low bin edge [GeV]	20	30	45	70	100	140	200	275	400	550
High bin edge [GeV]	30	45	70	100	140	200	275	400	550	1050

TABLE B.4: Differential cross-section measurements for EW- Zjj production and inclusive Zjj production as a function of $p_T^{\ell\ell}$. The EW- Zjj measurements are performed in the signal region. The inclusive Zjj measurements are performed in the signal region and three control regions. The different sources of uncertainty are grouped together for each measurement, with a more granular breakdown of each uncertainty available in HEPDATA.

Appendix C

EFT signal injection validation

This appendix describes a validation of the EFT interpretation results of section 7.3. The aim of this validation is to check that injection of the expected EFT signal from the 95% confidence intervals into Asimov data results in a proportional change in the extracted EW- Zjj signal.

Table 7.1 presents the observed 95% confidence limits from the likelihood fits for each of the four Wilson coefficients. These fits provide an upper and lower bound on the c_k , the impact on the EW- Zjj signal can be evaluated at these values according to Equation (7.3). Eight Asimov datasets were created, two for each of the four Wilson coefficients corresponding to the upper and lower limits of the observed confidence intervals. The strong- Zjj component in the Asimov data is from the MG5_NLO+PY8 prediction described in Section 6.3 (AsimovFxFx).

The EW signal was extracted from each of the modified Asimov datasets as described in Section 6.1 and compared to the EW signal extracted from the nominal Asimov dataset. The observed difference between extracted signals is compared to the injected signal for each dataset, with any deviation being interpreted as a bias of the method.

Wilson coefficient	Lower/Upper bound of 95 % CL	Injected EFT / Extracted EW	1	2	3	4	5	$\Delta\phi(j,j)$ bin		8	9	10	11	12
c_W	Lower	Injected	-13.0	-11.2	-9.4	20.0	18.1	-8.9	-8.4	17.7	19.9	-9.4	-10.7	-13.2
	Lower	Extracted	-9.4	-8.9	-7.7	20.0	16.9	-9.9	-9.4	16.6	19.9	-7.7	-8.4	-9.7
	Upper	Injected	30.6	26.8	26.7	-27.8	-23.6	34.0	32.6	-23.1	-27.8	26.9	25.7	31.0
	Upper	Extracted	27.3	23.8	22.4	-33.3	-26.4	32.4	30.9	-26.1	-33.4	22.5	22.7	27.8
$\tilde{c}_W$	Lower	Injected	-2.1	-6.1	-17.9	-21.9	12.2	11.3	-10.2	-10.7	23.1	18.8	6.1	2.4
	Lower	Extracted	-0.6	-5.1	-16.9	-21.0	12.7	11.5	-10.1	-10.2	24.5	20.4	7.5	4.1
	Upper	Injected	2.9	8.1	23.7	30.1	-13.2	-12.3	15.0	15.8	-27.2	-23.0	-7.4	-2.7
	Upper	Extracted	1.7	7.2	22.3	28.2	-14.5	-13.1	14.4	14.5	-29.7	-25.1	-8.9	-4.3
c_{HWB}	Lower	Injected	-50.7	-46.3	-49.9	-35.4	25.9	58.5	60.8	24.6	-31.0	-47.7	-35.1	-45.1
	Lower	Extracted	-59.2	-52.3	-58.1	-45.8	21.7	56.8	59.2	19.8	-42.2	-55.0	-40.7	-54.3
	Upper	Injected	56.3	47.8	54.5	44.0	9.7	-2.2	-2.9	10.0	43.3	54.8	45.1	55.4
	Upper	Extracted	56.8	48.2	55.5	45.2	10.1	-2.0	-2.6	10.6	44.5	55.6	45.5	56.0
$\tilde{c}_{\text{HWB}}$	Lower	Injected	-0.2	-0.5	-1.7	-5.0	-4.2	-1.5	1.5	4.1	4.9	1.7	0.5	0.2
	Lower	Extracted	-0.1	-0.4	-1.6	-4.9	-4.2	-1.5	1.5	4.1	5.0	1.8	0.5	0.3
	Upper	Injected	-1.7	-5.0	-17.2	-50.1	-41.2	-13.7	17.3	43.8	51.4	17.9	4.9	1.6
	Upper	Extracted	-0.7	-4.4	-16.5	-49.2	-41.1	-13.6	17.2	43.2	51.5	18.4	5.5	2.7

TABLE C.1: Injected EFT signal component and extracted EW signal modification as a function of $\Delta\phi_{jj}^{\text{signed}}$ corresponding to the upper and lower bounds of the extracted 95% confidence intervals measured for each of the four Wilson coefficients, extracted using AsimovFxFx. The event yields N_{SM} in each of the bins are (519, 419, 482, 379, 136, 76, 76, 136, 379, 482, 419, 519).

The injected EFT component per bin of $\Delta\phi_{jj}^{\text{signed}}$ and extracted difference in EW signal compared to the nominal Asimov fit are presented in Table C.1. The *injected* values presented corresponds to the $N_{\text{EFT}} - N_{\text{SM}}$, *i.e.* the difference in event yield introduced by the EFT variation and the *extracted* value is the difference in the number of EW events extracted with respect to the nominal Asimov fit. The EW extraction successfully extracts the EFT modifications in all cases. The only operator that shows a minor bias is c_W , this is due to the impact of this operator has a very non linear shape as can be seen in Figure 7.1 and the linear transfer function that is fit in the EW extraction is not flexible enough to precisely extract this EFT contribution to the signal.

Bibliography

- [1] The ATLAS Collaboration.
 “The ATLAS Experiment at the CERN Large Hadron Collider”.
 In: *JINST* 3 (2008), S08003. DOI: [10.1088/1748-0221/3/08/S08003](https://doi.org/10.1088/1748-0221/3/08/S08003).
- [2] The CMS Collaboration. “The CMS Experiment at the CERN LHC”.
 In: *JINST* 3 (2008), S08004. DOI: [10.1088/1748-0221/3/08/S08004](https://doi.org/10.1088/1748-0221/3/08/S08004).
- [3] G. Bachy et al. “The LEP collider: construction, project status and outlook”.
 In: *Part. Accel.* 26 (1990). Ed. by Yoshitaka Kimura et al., pp. 19–32.
- [4] Paolo Spagnolo. “Boson gauge couplings at LEP”.
 In: *40th Rencontres de Moriond on Electroweak Interactions and Unified Theories*.
 June 2005, pp. 363–370. arXiv: [hep-ex/0506048](https://arxiv.org/abs/hep-ex/0506048).
- [5] The ATLAS Collaboration.
 “Measurement of the electroweak production of dijets in association with a
 Z-boson and distributions sensitive to vector boson fusion in proton-proton
 collisions at $\sqrt{s} = 8$ TeV using the ATLAS detector”.
 In: *JHEP* 04 (2014), p. 031. DOI: [10.1007/JHEP04\(2014\)031](https://doi.org/10.1007/JHEP04(2014)031).
 arXiv: [1401.7610](https://arxiv.org/abs/1401.7610) [[hep-ex](https://arxiv.org/archive/hep)].
- [6] The CMS Collaboration. “Measurement of electroweak production of two
 jets in association with a Z boson in proton-proton collisions at $\sqrt{s} = 8$ TeV”.

- In: *Eur. Phys. J. C* 75.2 (2015), p. 66. DOI: [10.1140/epjc/s10052-014-3232-5](https://doi.org/10.1140/epjc/s10052-014-3232-5).
arXiv: [1410.3153](https://arxiv.org/abs/1410.3153) [hep-ex].
- [7] The ATLAS Collaboration. “Measurement of the cross-section for electroweak production of dijets in association with a Z boson in pp collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector”.
In: *Phys. Lett. B* 775 (2017), pp. 206–228.
DOI: [10.1016/j.physletb.2017.10.040](https://doi.org/10.1016/j.physletb.2017.10.040). arXiv: [1709.10264](https://arxiv.org/abs/1709.10264) [hep-ex].
- [8] The CMS Collaboration. “Electroweak production of two jets in association with a Z boson in proton–proton collisions at $\sqrt{s} = 13$ TeV”.
In: *Eur. Phys. J. C* 78.7 (2018), p. 589.
DOI: [10.1140/epjc/s10052-018-6049-9](https://doi.org/10.1140/epjc/s10052-018-6049-9). arXiv: [1712.09814](https://arxiv.org/abs/1712.09814) [hep-ex].
- [9] C.F. Anders et al.
“Vector boson scattering: Recent experimental and theory developments”.
In: *Rev. Phys.* 3 (2018), pp. 44–63. DOI: [10.1016/j.revip.2018.11.001](https://doi.org/10.1016/j.revip.2018.11.001).
arXiv: [1801.04203](https://arxiv.org/abs/1801.04203) [hep-ph].
- [10] Guido Altarelli. *Collider Physics within the Standard Model: A Primer*. Springer Nature, 2017.
- [11] David Griffiths. *Introduction to Elementary Particles*. Wiley, 1987.
- [12] Michael E. Peskin and Daniel V. Schroeder.
An Introduction to Quantum Field Theory. Westview Press, 1995.
- [13] Mark Thomson. *Modern Particle Physics*. Cambridge University Press, 2013.

- [14] J. H. Christenson et al. “Evidence for the 2π Decay of the K_2^0 Meson”.
In: *Phys. Rev. Lett.* 13 (4 1964), pp. 138–140.
DOI: [10.1103/PhysRevLett.13.138](https://doi.org/10.1103/PhysRevLett.13.138).
URL: <https://link.aps.org/doi/10.1103/PhysRevLett.13.138>.
- [15] C. N. Yang and R. L. Mills.
“Conservation of Isotopic Spin and Isotopic Gauge Invariance”.
In: *Phys. Rev.* 96 (1 1954), pp. 191–195. DOI: [10.1103/PhysRev.96.191](https://doi.org/10.1103/PhysRev.96.191).
URL: <https://link.aps.org/doi/10.1103/PhysRev.96.191>.
- [16] S.L. Glashow and Murray Gell-Mann.
“GAUGE THEORIES OF VECTOR PARTICLES”.
In: *Annals Phys.* 15 (1961), pp. 437–460.
DOI: [10.1016/0003-4916\(61\)90193-2](https://doi.org/10.1016/0003-4916(61)90193-2).
- [17] U. Baur and D. Zeppenfeld.
“Measuring three vector boson couplings in $q q \rightarrow q q W$ at the SSC”.
In: *Workshop on Physics at Current Accelerators and the Supercollider*. Sept. 1993,
pp. 0327–334. arXiv: [hep-ph/9309227](https://arxiv.org/abs/hep-ph/9309227).
- [18] M. Tanabashi et al. “Review of Particle Physics”.
In: *Phys. Rev. D* 98 (3 2018), p. 030001. DOI: [10.1103/PhysRevD.98.030001](https://doi.org/10.1103/PhysRevD.98.030001).
URL: <https://link.aps.org/doi/10.1103/PhysRevD.98.030001>.
- [19] Particle Data Group et al. “Review of Particle Physics”.
In: *Progress of Theoretical and Experimental Physics* 2020.8 (Aug. 2020). 083C01.
ISSN: 2050-3911. DOI: [10.1093/ptep/ptaa104](https://doi.org/10.1093/ptep/ptaa104).
eprint: <https://academic.oup.com/ptep/article->

- [pdf/2020/8/083C01/34673722/ptaa104.pdf](#).
URL: <https://doi.org/10.1093/ptep/ptaa104>.
- [20] *Standard Model Summary Plots Spring 2019*. Tech. rep. Geneva: CERN, 2019.
URL: <http://cds.cern.ch/record/2668559>.
- [21] L.A. Harland-Lang et al.
“Parton distributions in the LHC era: MMHT 2014 PDFs”.
In: *Eur. Phys. J. C* 75.5 (2015), p. 204.
DOI: [10.1140/epjc/s10052-015-3397-6](https://doi.org/10.1140/epjc/s10052-015-3397-6). arXiv: [1412.3989](https://arxiv.org/abs/1412.3989) [hep-ph].
- [22] William A. Bardeen et al. “Deep-inelastic scattering beyond the leading order in asymptotically free gauge theories”.
In: *Phys. Rev. D* 18 (11 1978), pp. 3998–4017.
DOI: [10.1103/PhysRevD.18.3998](https://doi.org/10.1103/PhysRevD.18.3998).
URL: <https://link.aps.org/doi/10.1103/PhysRevD.18.3998>.
- [23] Andy Buckley et al. “General-purpose event generators for LHC physics”.
In: *Phys. Rept.* 504 (2011), pp. 145–233.
DOI: [10.1016/j.physrep.2011.03.005](https://doi.org/10.1016/j.physrep.2011.03.005). arXiv: [1101.2599](https://arxiv.org/abs/1101.2599) [hep-ph].
- [24] S. Catani et al. “QCD matrix elements + parton showers”.
In: *JHEP* 11 (2001), p. 063. DOI: [10.1088/1126-6708/2001/11/063](https://doi.org/10.1088/1126-6708/2001/11/063).
arXiv: [hep-ph/0109231](https://arxiv.org/abs/hep-ph/0109231).
- [25] Stefan Hoeche et al. “QCD matrix elements and truncated showers”.
In: *JHEP* 05 (2009), p. 053. DOI: [10.1088/1126-6708/2009/05/053](https://doi.org/10.1088/1126-6708/2009/05/053).
arXiv: [0903.1219](https://arxiv.org/abs/0903.1219) [hep-ph].

- [26] Leif Lonnblad. “Correcting the color dipole cascade model with fixed order matrix elements”. In: *JHEP* 05 (2002), p. 046.
DOI: [10.1088/1126-6708/2002/05/046](https://doi.org/10.1088/1126-6708/2002/05/046). arXiv: [hep-ph/0112284](https://arxiv.org/abs/hep-ph/0112284).
- [27] Leif Lonnblad and Stefan Prestel.
“Matching Tree-Level Matrix Elements with Interleaved Showers”.
In: *JHEP* 03 (2012), p. 019. DOI: [10.1007/JHEP03\(2012\)019](https://doi.org/10.1007/JHEP03(2012)019).
arXiv: [1109.4829](https://arxiv.org/abs/1109.4829) [hep-ph].
- [28] Rikkert Frederix and Stefano Frixione.
“Merging meets matching in MC@NLO”. In: *JHEP* 12 (2012), p. 061.
DOI: [10.1007/JHEP12\(2012\)061](https://doi.org/10.1007/JHEP12(2012)061). arXiv: [1209.6215](https://arxiv.org/abs/1209.6215) [hep-ph].
- [29] Bo Andersson. “The Lund model”.
In: *Camb. Monogr. Part. Phys. Nucl. Phys. Cosmol.* 7 (1997), pp. 1–471.
- [30] B. Andersson et al. “Parton fragmentation and string dynamics”.
In: *Physics Reports* 97.2 (1983), pp. 31 –145. ISSN: 0370-1573.
DOI: [https://doi.org/10.1016/0370-1573\(83\)90080-7](https://doi.org/10.1016/0370-1573(83)90080-7). URL: <http://www.sciencedirect.com/science/article/pii/0370157383900807>.
- [31] B.R. Webber.
“A QCD Model for Jet Fragmentation Including Soft Gluon Interference”.
In: *Nucl. Phys. B* 238 (1984), pp. 492–528.
DOI: [10.1016/0550-3213\(84\)90333-X](https://doi.org/10.1016/0550-3213(84)90333-X).
- [32] Xianguo Lu. “Measurement of hadron composition in charged jets from pp collisions with the ALICE experiment”.
In: *Nuclear Physics A* 931 (2014). QUARK MATTER 2014, pp. 428 –432.

ISSN: 0375-9474.

DOI: <https://doi.org/10.1016/j.nuclphysa.2014.08.003>. URL: <http://www.sciencedirect.com/science/article/pii/S0375947414002528>.

- [33] T. Gleisberg et al. “Event generation with SHERPA 1.1”.
In: *JHEP* 02 (2009), p. 007. DOI: [10.1088/1126-6708/2009/02/007](https://doi.org/10.1088/1126-6708/2009/02/007).
arXiv: [0811.4622](https://arxiv.org/abs/0811.4622) [hep-ph].
- [34] The ATLAS collaboration. *Combined Muon Performance Public Results*. 2020.
URL: <https://twiki.cern.ch/twiki/bin/view/AtlasPublic/MuonPerformancePublicPlotsOld> (visited on 08/17/2020).
- [35] The ATLAS Collaboration.
“Measurements of the production cross section of a Z boson in association with jets in pp collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector”.
In: *Eur. Phys. J. C* 77.6 (2017), p. 361.
DOI: [10.1140/epjc/s10052-017-4900-z](https://doi.org/10.1140/epjc/s10052-017-4900-z). arXiv: [1702.05725](https://arxiv.org/abs/1702.05725) [hep-ex].
- [36] Carlo Oleari and Dieter Zeppenfeld.
“QCD corrections to electroweak $l\nu l j$ and $l+l-jj$ production”.
In: *Physical Review D* 69.9 (2004). ISSN: 1550-2368.
DOI: [10.1103/physrevd.69.093004](https://doi.org/10.1103/physrevd.69.093004).
URL: <http://dx.doi.org/10.1103/PhysRevD.69.093004>.
- [37] A. Dedes et al.
“Feynman rules for the Standard Model Effective Field Theory in R -gauges”.
In: *JHEP* 06 (2017), p. 143. DOI: [10.1007/JHEP06\(2017\)143](https://doi.org/10.1007/JHEP06(2017)143).
arXiv: [1704.03888](https://arxiv.org/abs/1704.03888) [hep-ph].

- [38] Raquel Gomez-Ambrosio.
“Studies of Dimension-Six EFT effects in Vector Boson Scattering”.
In: *Eur. Phys. J. C* 79.5 (2019), p. 389.
DOI: [10.1140/epjc/s10052-019-6893-2](https://doi.org/10.1140/epjc/s10052-019-6893-2). arXiv: [1809.04189](https://arxiv.org/abs/1809.04189) [hep-ph].
- [39] Oliver Sim Brüning et al. *LHC Design Report*.
CERN Yellow Reports: Monographs. Geneva: CERN, 2004.
DOI: [10.5170/CERN-2004-003-V-1](https://doi.org/10.5170/CERN-2004-003-V-1).
URL: <https://cds.cern.ch/record/782076>.
- [40] Lyndon Evans and Philip Bryant. “LHC Machine”.
In: *Journal of Instrumentation* 3.08 (2008), S08001–S08001.
DOI: [10.1088/1748-0221/3/08/s08001](https://doi.org/10.1088/1748-0221/3/08/s08001).
URL: <https://doi.org/10.1088/1748-0221/3/08/s08001>.
- [41] Esma Mobs.
“The CERN accelerator complex. Complexe des accélérateurs du CERN”.
In: (2016). General Photo. URL: <https://cds.cern.ch/record/2197559>.
- [42] The LHCb Collaboration. “The LHCb Detector at the LHC”.
In: *JINST* 3 (2008), S08005. DOI: [10.1088/1748-0221/3/08/S08005](https://doi.org/10.1088/1748-0221/3/08/S08005).
- [43] The ALICE Collaboration. “The ALICE experiment at the CERN LHC”.
In: *JINST* 3 (2008), S08002. DOI: [10.1088/1748-0221/3/08/S08002](https://doi.org/10.1088/1748-0221/3/08/S08002).
- [44] by Achintya Rao and Achintya Rao.
“The 2018 data-taking run at the LHC has begun”. In: (2018).
URL: <http://cds.cern.ch/record/2315786>.

- [45] A. Ruiz Martínez. “The Run-2 ATLAS Trigger System”. In: *J. Phys. Conf. Ser.* 762.1 (2016). Ed. by Luis Salinas and Claudio Torres, p. 012003.
DOI: [10.1088/1742-6596/762/1/012003](https://doi.org/10.1088/1742-6596/762/1/012003).
- [46] Martin zur Nedden.
The Run-2 ATLAS Trigger System: Design, Performance and Plan.
Tech. rep. ATL-DAQ-PROC-2016-039. Geneva: CERN, 2016.
URL: <https://cds.cern.ch/record/2238679>.
- [47] G. Avoni et al. “The new LUCID-2 detector for luminosity measurement and monitoring in ATLAS”.
In: *Journal of Instrumentation* 13.07 (2018), P07017–P07017.
DOI: [10.1088/1748-0221/13/07/p07017](https://doi.org/10.1088/1748-0221/13/07/p07017).
URL: <https://doi.org/10.1088/1748-0221/13/07/p07017>.
- [48] *Luminosity determination in pp collisions at $\sqrt{s} = 13$ TeV using the ATLAS detector at the LHC*. Tech. rep. ATLAS-CONF-2019-021. Geneva: CERN, 2019.
URL: <http://cds.cern.ch/record/2677054>.
- [49] The ATLAS collaboration. *Luminosity Public Results Run 2*. 2020.
URL: <https://twiki.cern.ch/twiki/bin/view/AtlasPublic/LuminosityPublicResultsRun2> (visited on 07/14/2020).
- [50] *Track Reconstruction Performance of the ATLAS Inner Detector at $\sqrt{s} = 13$ TeV*.
Tech. rep. ATL-PHYS-PUB-2015-018. Geneva: CERN, 2015.
URL: <http://cds.cern.ch/record/2037683>.
- [51] Norbert Wermes and G Hallelwel.
ATLAS pixel detector: Technical Design Report.

- Technical Design Report ATLAS. Geneva: CERN, 1998.
URL: <https://cds.cern.ch/record/381263>.
- [52] M Capeans et al. *ATLAS Insertable B-Layer Technical Design Report*.
Tech. rep. CERN-LHCC-2010-013. ATLAS-TDR-19. 2010.
URL: <https://cds.cern.ch/record/1291633>.
- [53] The ATLAS Collaboration.
“Operation and performance of the ATLAS semiconductor tracker”.
In: *JINST* 9 (2014), P08009. DOI: [10.1088/1748-0221/9/08/P08009](https://doi.org/10.1088/1748-0221/9/08/P08009).
arXiv: [1404.7473](https://arxiv.org/abs/1404.7473) [hep-ex].
- [54] The ATLAS Collaboration. “The ATLAS Transition Radiation Tracker (TRT) proportional drift tube: Design and performance”.
In: *JINST* 3 (2008), P02013. DOI: [10.1088/1748-0221/3/02/P02013](https://doi.org/10.1088/1748-0221/3/02/P02013).
- [55] *ATLAS liquid-argon calorimeter: Technical Design Report*.
Technical Design Report ATLAS. Geneva: CERN, 1996.
URL: <https://cds.cern.ch/record/331061>.
- [56] *ATLAS tile calorimeter: Technical Design Report*.
Technical Design Report ATLAS. Geneva: CERN, 1996.
URL: <https://cds.cern.ch/record/331062>.
- [57] *ATLAS muon spectrometer: Technical Design Report*.
Technical Design Report ATLAS. Geneva: CERN, 1997.
URL: <https://cds.cern.ch/record/331068>.
- [58] C. Ferretti and H. Kroha.
“Upgrades of the ATLAS Muon Spectrometer with sMDT Chambers”.

- In: *Nucl. Instrum. Meth. A* 824 (2016). Ed. by Maria Giuseppina Bisogni et al., pp. 538–540. DOI: [10.1016/j.nima.2015.11.043](https://doi.org/10.1016/j.nima.2015.11.043).
arXiv: [1603.09544](https://arxiv.org/abs/1603.09544) [physics.ins-det].
- [59] The ATLAS Collaboration. “Optimisation of large-radius jet reconstruction for the ATLAS detector in 13 TeV proton–proton collisions”. In: (Sept. 2020).
arXiv: [2009.04986](https://arxiv.org/abs/2009.04986) [hep-ex].
- [60] “Luminosity determination in pp collisions at $\sqrt{s} = 13$ TeV using the ATLAS detector at the LHC”. In: (June 2019).
- [61] The ATLAS collaboration. *How To Clean Jets 2017*. 2020. URL: <https://twiki.cern.ch/twiki/bin/view/AtlasProtected/HowToCleanJets2017> (visited on 07/28/2020).
- [62] Paolo Nason. “A New method for combining NLO QCD with shower Monte Carlo algorithms”. In: *JHEP* 11 (2004), p. 040.
DOI: [10.1088/1126-6708/2004/11/040](https://doi.org/10.1088/1126-6708/2004/11/040). arXiv: [hep-ph/0409146](https://arxiv.org/abs/hep-ph/0409146).
- [63] Stefano Frixione, Paolo Nason, and Carlo Oleari. “Matching NLO QCD computations with Parton Shower simulations: the POWHEG method”. In: *JHEP* 11 (2007), p. 070. DOI: [10.1088/1126-6708/2007/11/070](https://doi.org/10.1088/1126-6708/2007/11/070).
arXiv: [0709.2092](https://arxiv.org/abs/0709.2092) [hep-ph].
- [64] Simone Alioli et al. “A general framework for implementing NLO calculations in shower Monte Carlo programs: the POWHEG BOX”. In: *JHEP* 06 (2010), p. 043. DOI: [10.1007/JHEP06\(2010\)043](https://doi.org/10.1007/JHEP06(2010)043).
arXiv: [1002.2581](https://arxiv.org/abs/1002.2581) [hep-ph].

- [65] J. Alwall et al. “The automated computation of tree-level and next-to-leading order differential cross sections, and their matching to parton shower simulations”. In: *JHEP* 07 (2014), p. 079.
DOI: [10.1007/JHEP07\(2014\)079](https://doi.org/10.1007/JHEP07(2014)079). arXiv: [1405.0301](https://arxiv.org/abs/1405.0301) [hep-ph].
- [66] Barbara Jager, Steven Schneider, and Giulia Zanderighi.
“Next-to-leading order QCD corrections to electroweak Zjj production in the POWHEG BOX”. In: *JHEP* 09 (2012), p. 083.
DOI: [10.1007/JHEP09\(2012\)083](https://doi.org/10.1007/JHEP09(2012)083). arXiv: [1207.2626](https://arxiv.org/abs/1207.2626) [hep-ph].
- [67] The ATLAS Collaboration.
“Measurement of the Z/γ^* boson transverse momentum distribution in pp collisions at $\sqrt{s} = 7$ TeV with the ATLAS detector”. In: *JHEP* 09 (2014), p. 145.
DOI: [10.1007/JHEP09\(2014\)145](https://doi.org/10.1007/JHEP09(2014)145). arXiv: [1406.3660](https://arxiv.org/abs/1406.3660) [hep-ex].
- [68] David J. Lange. “The EvtGen particle decay simulation package”.
In: *Nuclear Instruments and Methods in Physics Research Section A: Accelerators, Spectrometers, Detectors and Associated Equipment* 462.1 (2001). BEAUTY2000, Proceedings of the 7th Int. Conf. on B-Physics at Hadron Machines, pp. 152–155. ISSN: 0168-9002.
DOI: [https://doi.org/10.1016/S0168-9002\(01\)00089-4](https://doi.org/10.1016/S0168-9002(01)00089-4). URL: <http://www.sciencedirect.com/science/article/pii/S0168900201000894>.
- [69] M. Bahr et al. “Herwig++ Physics and Manual”.
In: *Eur. Phys. J. C* 58 (2008), pp. 639–707.
DOI: [10.1140/epjc/s10052-008-0798-9](https://doi.org/10.1140/epjc/s10052-008-0798-9). arXiv: [0803.0883](https://arxiv.org/abs/0803.0883) [hep-ph].

- [70] Johannes Bellm et al. “Herwig 7.0/Herwig++ 3.0 release note”.
In: *Eur. Phys. J. C* 76.4 (2016), p. 196.
DOI: [10.1140/epjc/s10052-016-4018-8](https://doi.org/10.1140/epjc/s10052-016-4018-8). arXiv: [1512.01178](https://arxiv.org/abs/1512.01178) [hep-ph].
- [71] J. Baglio et al. “VBFNLO: A Parton Level Monte Carlo for Processes with Electroweak Bosons – Manual for Version 2.7.0”. In: (July 2011).
arXiv: [1107.4038](https://arxiv.org/abs/1107.4038) [hep-ph].
- [72] Enrico Bothmann et al. “Event Generation with Sherpa 2.2”.
In: *SciPost Phys.* 7.3 (2019), p. 034. DOI: [10.21468/SciPostPhys.7.3.034](https://doi.org/10.21468/SciPostPhys.7.3.034).
arXiv: [1905.09127](https://arxiv.org/abs/1905.09127) [hep-ph].
- [73] Richard D. Ball et al. “Parton distributions for the LHC Run II”.
In: *JHEP* 04 (2015), p. 040. DOI: [10.1007/JHEP04\(2015\)040](https://doi.org/10.1007/JHEP04(2015)040).
arXiv: [1410.8849](https://arxiv.org/abs/1410.8849) [hep-ph].
- [74] *Modelling of the vector boson scattering process $pp \rightarrow W^\pm W^\pm jj$ in Monte Carlo generators in ATLAS*. Tech. rep. ATL-PHYS-PUB-2019-004.
Geneva: CERN, 2019. URL: <https://cds.cern.ch/record/2655303>.
- [75] Tanju Gleisberg and Stefan Hoeche.
“Comix, a new matrix element generator”. In: *JHEP* 12 (2008), p. 039.
DOI: [10.1088/1126-6708/2008/12/039](https://doi.org/10.1088/1126-6708/2008/12/039). arXiv: [0808.3674](https://arxiv.org/abs/0808.3674) [hep-ph].
- [76] Ansgar Denner, Stefan Dittmaier, and Lars Hofer. “Collier: a fortran-based Complex One-Loop Library in Extended Regularizations”.
In: *Comput. Phys. Commun.* 212 (2017), pp. 220–238.
DOI: [10.1016/j.cpc.2016.10.013](https://doi.org/10.1016/j.cpc.2016.10.013). arXiv: [1604.06792](https://arxiv.org/abs/1604.06792) [hep-ph].

-
- [77] Fabio Cascioli, Philipp Maierhofer, and Stefano Pozzorini.
“Scattering Amplitudes with Open Loops”.
In: *Phys. Rev. Lett.* 108 (2012), p. 111601.
DOI: [10.1103/PhysRevLett.108.111601](https://doi.org/10.1103/PhysRevLett.108.111601). arXiv: [1111.5206 \[hep-ph\]](https://arxiv.org/abs/1111.5206).
- [78] Stefan Hoeche et al.
“QCD matrix elements + parton showers: The NLO case”.
In: *JHEP* 04 (2013), p. 027. DOI: [10.1007/JHEP04\(2013\)027](https://doi.org/10.1007/JHEP04(2013)027).
arXiv: [1207.5030 \[hep-ph\]](https://arxiv.org/abs/1207.5030).
- [79] Charalampos Anastasiou et al. “High precision QCD at hadron colliders: Electroweak gauge boson rapidity distributions at NNLO”.
In: *Phys. Rev. D* 69 (2004), p. 094008. DOI: [10.1103/PhysRevD.69.094008](https://doi.org/10.1103/PhysRevD.69.094008).
arXiv: [hep-ph/0312266](https://arxiv.org/abs/hep-ph/0312266).
- [80] Richard D. Ball et al. “Parton distributions with LHC data”.
In: *Nucl. Phys. B* 867 (2013), pp. 244–289.
DOI: [10.1016/j.nuclphysb.2012.10.003](https://doi.org/10.1016/j.nuclphysb.2012.10.003). arXiv: [1207.1303 \[hep-ph\]](https://arxiv.org/abs/1207.1303).
- [81] *ATLAS Pythia 8 tunes to 7 TeV datas*. Tech. rep. ATL-PHYS-PUB-2014-021.
Geneva: CERN, 2014. URL: <https://cds.cern.ch/record/1966419>.
- [82] Stefano Frixione, Paolo Nason, and Giovanni Ridolfi. “A Positive-weight next-to-leading-order Monte Carlo for heavy flavour hadroproduction”.
In: *JHEP* 09 (2007), p. 126. DOI: [10.1088/1126-6708/2007/09/126](https://doi.org/10.1088/1126-6708/2007/09/126).
arXiv: [0707.3088 \[hep-ph\]](https://arxiv.org/abs/0707.3088).
- [83] S. Agostinelli et al. “Geant4—a simulation toolkit”.
In: *Nuclear Instruments and Methods in Physics Research Section A: Accelerators,*

- Spectrometers, Detectors and Associated Equipment* 506.3 (2003), pp. 250–303.
ISSN: 0168-9002. DOI: [https://doi.org/10.1016/S0168-9002\(03\)01368-8](https://doi.org/10.1016/S0168-9002(03)01368-8).
URL: <http://www.sciencedirect.com/science/article/pii/S0168900203013688>.
- [84] The ATLAS Collaboration. “The ATLAS Simulation Infrastructure”.
In: *Eur. Phys. J. C* 70 (2010), pp. 823–874.
DOI: [10.1140/epjc/s10052-010-1429-9](https://doi.org/10.1140/epjc/s10052-010-1429-9).
arXiv: [1005.4568](https://arxiv.org/abs/1005.4568) [physics.ins-det].
- [85] Torbjorn Sjostrand, Stephen Mrenna, and Peter Z. Skands.
“A Brief Introduction to PYTHIA 8.1”.
In: *Comput. Phys. Commun.* 178 (2008), pp. 852–867.
DOI: [10.1016/j.cpc.2008.01.036](https://doi.org/10.1016/j.cpc.2008.01.036). arXiv: [0710.3820](https://arxiv.org/abs/0710.3820) [hep-ph].
- [86] *The Pythia 8 A3 tune description of ATLAS minimum bias and inelastic measurements incorporating the Donnachie-Landshoff diffractive model*.
Tech. rep. ATL-PHYS-PUB-2016-017. Geneva: CERN, 2016.
URL: <https://cds.cern.ch/record/2206965>.
- [87] The ATLAS collaboration. *Event Displays from Run 2 physics analyses*. 2020.
URL: <https://twiki.cern.ch/twiki/bin/view/AtlasPublic/EventDisplayRun2Physics> (visited on 08/28/2020).
- [88] R. Fruhwirth. “Application of Kalman filtering to track and vertex fitting”.
In: *Nucl. Instrum. Meth. A* 262 (1987), pp. 444–450.
DOI: [10.1016/0168-9002\(87\)90887-4](https://doi.org/10.1016/0168-9002(87)90887-4).

- [89] *Performance of the ATLAS Inner Detector Track and Vertex Reconstruction in the High Pile-Up LHC Environment*. Tech. rep. ATLAS-CONF-2012-042. Geneva: CERN, 2012. URL: <https://cds.cern.ch/record/1435196>.
- [90] The ATLAS collaboration. *Tracking CP Recommendations*. 2020. URL: <https://twiki.cern.ch/twiki/bin/view/AtlasProtected/TrackingCPRecsEarly2018> (visited on 08/28/2020).
- [91] The ATLAS Collaboration. “Reconstruction of primary vertices at the ATLAS experiment in Run 1 proton–proton collisions at the LHC”. In: *Eur. Phys. J. C* 77.5 (2017), p. 332. DOI: [10.1140/epjc/s10052-017-4887-5](https://doi.org/10.1140/epjc/s10052-017-4887-5). arXiv: [1611.10235](https://arxiv.org/abs/1611.10235) [physics.ins-det].
- [92] The ATLAS Collaboration. “Electron and photon performance measurements with the ATLAS detector using the 2015–2017 LHC proton-proton collision data”. In: *JINST* 14.12 (2019), P12006. DOI: [10.1088/1748-0221/14/12/P12006](https://doi.org/10.1088/1748-0221/14/12/P12006). arXiv: [1908.00005](https://arxiv.org/abs/1908.00005) [hep-ex].
- [93] The ATLAS Collaboration. “Muon reconstruction performance of the ATLAS detector in proton–proton collision data at $\sqrt{s}=13$ TeV”. In: *Eur. Phys. J. C* 76.5 (2016), p. 292. DOI: [10.1140/epjc/s10052-016-4120-y](https://doi.org/10.1140/epjc/s10052-016-4120-y). arXiv: [1603.05598](https://arxiv.org/abs/1603.05598) [hep-ex].
- [94] Matteo Cacciari, Gavin P. Salam, and Gregory Soyez. “The anti- k_t jet clustering algorithm”. In: *JHEP* 04 (2008), p. 063. DOI: [10.1088/1126-6708/2008/04/063](https://doi.org/10.1088/1126-6708/2008/04/063). arXiv: [0802.1189](https://arxiv.org/abs/0802.1189) [hep-ph].

- [95] Matteo Cacciari, Gavin P. Salam, and Gregory Soyez. “FastJet User Manual”. In: *Eur. Phys. J. C* 72 (2012), p. 1896.
DOI: [10.1140/epjc/s10052-012-1896-2](https://doi.org/10.1140/epjc/s10052-012-1896-2). arXiv: [1111.6097](https://arxiv.org/abs/1111.6097) [hep-ph].
- [96] S. Catani et al. “Longitudinally invariant K_t clustering algorithms for hadron hadron collisions”. In: *Nucl. Phys. B* 406 (1993), pp. 187–224.
DOI: [10.1016/0550-3213\(93\)90166-M](https://doi.org/10.1016/0550-3213(93)90166-M).
- [97] Stephen D. Ellis and Davison E. Soper.
“Successive combination jet algorithm for hadron collisions”.
In: *Phys. Rev. D* 48 (1993), pp. 3160–3166. DOI: [10.1103/PhysRevD.48.3160](https://doi.org/10.1103/PhysRevD.48.3160).
arXiv: [hep-ph/9305266](https://arxiv.org/abs/hep-ph/9305266).
- [98] The ATLAS Collaboration. “Topological cell clustering in the ATLAS calorimeters and its performance in LHC Run 1”.
In: *Eur. Phys. J. C* 77 (2017), p. 490. DOI: [10.1140/epjc/s10052-017-5004-5](https://doi.org/10.1140/epjc/s10052-017-5004-5).
arXiv: [1603.02934](https://arxiv.org/abs/1603.02934) [hep-ex].
- [99] The ATLAS Collaboration. “Jet reconstruction and performance using particle flow with the ATLAS Detector”. In: *Eur. Phys. J. C* 77.7 (2017), p. 466.
DOI: [10.1140/epjc/s10052-017-5031-2](https://doi.org/10.1140/epjc/s10052-017-5031-2). arXiv: [1703.10485](https://arxiv.org/abs/1703.10485) [hep-ex].
- [100] The ATLAS Collaboration.
“Jet energy scale measurements and their systematic uncertainties in proton-proton collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector”.
In: *Phys. Rev. D* 96.7 (2017), p. 072002. DOI: [10.1103/PhysRevD.96.072002](https://doi.org/10.1103/PhysRevD.96.072002).
arXiv: [1703.09665](https://arxiv.org/abs/1703.09665) [hep-ex].

-
- [101] *Tagging and suppression of pileup jets with the ATLAS detector*.
Tech. rep. ATLAS-CONF-2014-018. Geneva: CERN, 2014.
URL: <https://cds.cern.ch/record/1700870>.
- [102] Julia Gonski. *Jet Cleaning in 2016 and the Event Level Cleaning Tool*.
Tech. rep. ATL-COM-PHYS-2017-982. Geneva: CERN, 2017.
URL: <https://cds.cern.ch/record/2272136>.
- [103] William Buttinger. *Using Event Weights to account for differences in Instantaneous Luminosity and Trigger Prescale in Monte Carlo and Data*.
Tech. rep. ATL-COM-SOFT-2015-119. Geneva: CERN, 2015.
URL: <https://cds.cern.ch/record/2014726>.
- [104] V. Hankele et al. “Anomalous Higgs boson couplings in vector boson fusion at the CERN LHC”. In: *Physical Review D* 74.9 (2006). ISSN: 1550-2368.
DOI: [10.1103/physrevd.74.095001](https://doi.org/10.1103/physrevd.74.095001).
URL: <http://dx.doi.org/10.1103/PhysRevD.74.095001>.
- [105] The ATLAS Collaboration. “Measurements of electroweak Wjj production and constraints on anomalous gauge couplings with the ATLAS detector”. In: *Eur. Phys. J. C* 77.7 (2017), p. 474.
DOI: [10.1140/epjc/s10052-017-5007-2](https://doi.org/10.1140/epjc/s10052-017-5007-2). arXiv: [1703.04362](https://arxiv.org/abs/1703.04362) [hep-ex].
- [106] W. Buttinger. *TRooFit analysis package*. 2020.
URL: <https://gitlab.cern.ch/will/TRooFit> (visited on 09/28/2020).
- [107] Persi Diaconis and Bradley Efron.
“Computer-Intensive Methods in Statistics”.

- In: *Scientific American* 248.5 (1983), pp. 116–131. ISSN: 00368733, 19467087.
URL: <http://www.jstor.org/stable/24968902>.
- [108] The ATLAS collaboration. *Extended pileup reweighting*. 2020.
URL: <https://twiki.cern.ch/twiki/bin/viewauth/AtlasProtected/ExtendedPileupReweighting> (visited on 07/28/2020).
- [109] Christian Bierlich et al. “Robust Independent Validation of Experiment and Theory: Rivet version 3”. In: *SciPost Phys.* 8 (2020), p. 026.
DOI: [10.21468/SciPostPhys.8.2.026](https://doi.org/10.21468/SciPostPhys.8.2.026). arXiv: [1912.05451](https://arxiv.org/abs/1912.05451) [hep-ph].
- [110] G. D’Agostini.
“A multidimensional unfolding method based on Bayes’ theorem”.
In: *Nuclear Instruments and Methods in Physics Research Section A: Accelerators, Spectrometers, Detectors and Associated Equipment* 362.2 (1995), pp. 487–498.
ISSN: 0168-9002. DOI: [https://doi.org/10.1016/0168-9002\(95\)00274-X](https://doi.org/10.1016/0168-9002(95)00274-X).
URL: <http://www.sciencedirect.com/science/article/pii/016890029500274X>.
- [111] Tim Adye. “Unfolding algorithms and tests using RooUnfold”.
In: *PHYSTAT 2011*. Geneva: CERN, 2011, pp. 313–318.
DOI: [10.5170/CERN-2011-006.313](https://doi.org/10.5170/CERN-2011-006.313). arXiv: [1105.1160](https://arxiv.org/abs/1105.1160) [physics.data-an].
- [112] A Armbruster et al. *Practical considerations for unfolding*.
Tech. rep. ATL-COM-PHYS-2014-277. Geneva: CERN, 2014.
URL: <https://cds.cern.ch/record/1694351>.

-
- [113] ATLAS collaboration. *HEPDATA record of the analysis results*. 2020.
URL: <https://https://www.hepdata.net/record/ins1803608> (visited on 09/15/2020).
- [114] Ilaria Brivio, Yun Jiang, and Michael Trott.
“The SMEFTsim package, theory and tools”. In: *JHEP* 12 (2017), p. 070.
DOI: [10.1007/JHEP12\(2017\)070](https://doi.org/10.1007/JHEP12(2017)070). arXiv: [1709.06492](https://arxiv.org/abs/1709.06492) [hep-ph].
- [115] Gary J. Feldman and Robert D. Cousins.
“A Unified approach to the classical statistical analysis of small signals”.
In: *Phys. Rev. D* 57 (1998), pp. 3873–3889. DOI: [10.1103/PhysRevD.57.3873](https://doi.org/10.1103/PhysRevD.57.3873).
arXiv: [physics/9711021](https://arxiv.org/abs/hep-ph/9711021).
- [116] S. S. Wilks. “The Large-Sample Distribution of the Likelihood Ratio for Testing Composite Hypotheses”.
In: *Ann. Math. Statist.* 9.1 (Mar. 1938), pp. 60–62.
DOI: [10.1214/aoms/1177732360](https://doi.org/10.1214/aoms/1177732360).
URL: <https://doi.org/10.1214/aoms/1177732360>.