

THE SEARCH FOR EXOTIC DIBOSON PRODUCTION IN THE SEMILEPTONIC
CHANNELS WITH THE ATLAS DETECTOR

by

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Abstract

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This thesis presents several studies in the search for new physics in the production of electroweak gauge bosons pairs with 36 fb^{-1} of proton-proton collisions measured by the ATLAS detector. Processes with electroweak gauge bosons in the final state are sensitive to new physics which alter the electroweak or Higgs sector of the Standard Model. The studies are conducted in the “semileptonic” decay channels, where one of the electroweak gauge bosons decays hadronically and the other decays leptonically. The groundwork of these studies is established in a search for a new resonant particle which can decay on-shell to pairs of W/Z bosons in the $\ell\nu qq$ channel. No significant excess of data with respect to the background prediction was observed. Therefore, upper limits at the 95% CL_s confidence level are placed on the possible resonant masses in a strongly coupled Heavy Vector Triplet model and on Bulk Randall-Sundrum Gravitons at 3 TeV and 1.7 TeV, respectively. These results are expanded upon with a dedicated study on an ATLAS-wide combination of resonant diboson and dilepton search results. The combination improves the 95% CL_s upper excluded mass values to 5.5 TeV and 2.1 TeV, respectively. In conjunction, a search for non-resonant new physics is conducted through a measurement of the electroweak vector-boson scattering in all semileptonic channels. Evidence for the process was observed at a significance of 2.7σ and a fiducial cross-section measurement of $\sigma = 45.1 \pm 8.6 \text{ (stat.)}_{-14.6}^{+15.9} \text{ (syst.) fb}$ was extracted, consistent with the Standard Model prediction. Lastly, both the prospects of the $\ell\nu qq$ resonance search and the measurement of vector-boson scattering in the High-Luminosity LHC era were evaluated. With 3000 fb^{-1} of proton-proton data, the upper mass limits on new resonances in the $\ell\nu qq$ channel are expected to increase by 1.3 - 2 TeV depending on the benchmark model and the vector-boson scattering cross-section is expected to be measured at the percent level.

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Chapter 1

Introduction

A well-known scientist (some say it was Bertrand Russell) once gave a public lecture on astronomy. He described how the earth orbits around the sun and how the sun, in turn, orbits around the center of a vast collection of stars called our galaxy. At the end of the lecture, a little old lady at the back of the room got up and said: “What you have told us is rubbish. The world is really a flat plate supported on the back of a giant tortoise.” The scientist gave a superior smile before replying, “What is the tortoise standing on?” “You’re very clever, young man, very clever,” said the old lady. “But it’s turtles all the way down!”

—Stephen Hawking, *A Brief History of Time*

It became well understood in the mid-19th century that macroscopic matter is composed of constituents known as atoms [1, 2]. In 1897 J.J. Thompson discovered that cathode rays, actually composed of electrons, did not fall into this paradigm [3]. Not soon after in 1909, Rutherford discovered the existence of protons in the center of atoms [4, 5], and in 1932 Chadwick showed there were also neutral particles, known as neutrons, in the nuclei as well [6]. Subsequent experiments in deep inelastic scattering in the 1960’s and 1970’s further showed that the neutron and proton were also not fundamental particles; they were composed out of constituents known as quarks and gluons [7].

The field of particle physics attempts to study the end of this and other similar chains of observations. It is our understanding of nature at the smallest and most fundamental scale. In a reductionist sense, it is the “bleeding edge” where all larger-scale physical models should obey the laws set by it¹. Particle physics helps us explain aspects of nuclear science, cosmology, and the energetic interaction of light with matter (in X-ray or CAT scans for instance). To probe smaller length scales requires the energy of the probe to be larger. By studying physics at smaller scales, we are therefore also probing our understanding of nature just after the Big Bang when energy density of the universe was higher than what it is currently.

Our current understanding of the field is summarized by a theory known as the Standard

¹Importantly, macroscopic physics will probably not be practically calculable by particle physics, just implicitly obey its observations.

Model of particle physics. The Standard Model is a Quantum Field Theory, where forces and interactions are understood to originate from the exchange of particles. The Standard Model explains the observed matter content of the universe, alongside a description of the electromagnetic force, which governs a large fraction of macroscopic phenomena, the weak nuclear force, which describe radioactive decays, and the strong nuclear force which explains the formation of nuclei. With the discovery of the Higgs boson in 2012, all particles predicted by the Standard Model have been experimentally observed.

While the Standard Model has proven to be extremely fruitful in describing the results of almost all particle physics experiments, several observation in tension with the theory have been found. These include the observed amount of matter-antimatter asymmetry in the universe, and the lack of Dark Matter candidates. In addition, there are several theoretical shortcomings of the Standard Model, such as a failure to implement an understanding of gravity and a correlated issue known as the hierarchy problem, which involves the unexplained orders of magnitude difference between the scale of gravity and the rest of the Standard Model. These facts indicate that the Standard Model may not be a complete understanding of nature and some new physics is required. Theories which provide expansions or changes to the Standard Model are known as physics Beyond the Standard Model.

A complete experimental probe of physics Beyond the Standard Model is impossible as the nature of these models can vary vastly. It is then useful to focus studies on commonly occurring experimental signatures. Many new physics models attempt to address the issue of the hierarchy problem by expanding on the Higgs mechanism. Others expand on the nature of gravity or change the gauge structure of the Standard Model. These features in the Standard Model are intimately connected with the nature of electroweak gauge bosons, so alterations by new physics are expected to produce measurable differences in the production rate of processes with W/Z bosons.

The ATLAS experiment at the Large Hadron Collider (LHC) is currently one of the optimal experiments to search for signatures of this form. This is due to two features. The first is that the LHC provides the most energetic man-made particle collisions to date, which is important for probing Beyond the Standard Model physics which is less constrained at higher energies. The other is the collisions at the LHC occur at rates² in excess of 1 GHz, allowing for the study of rare processes.

This thesis will detail several studies on the search for new physics in the production of pairs of electroweak gauge bosons (dibosons) with the ATLAS detector using 13 TeV proton-proton data. The first several chapters will detail the theoretical and experimental background needed in this work. Chapter 2 will detail the relevant aspects of the Standard Model and several Beyond the Standard Model theories to be probed. Experimental details of the LHC and ATLAS detector are discussed in Chapter 3 and the techniques used to extract physics

²We will use SI units and prefixes throughout this work, except when discussing particle masses, energies, and momenta. These quantities will instead be provided in “Natural Units” where all three quantities are in units of electron-volt (eV).

information from the ATLAS data are discussed in Chapter 4.

The search for new physics in diboson final states with 36 fb^{-1} of 13 TeV ATLAS data is divided into several chapters. In Chapter 5 a search is conducted for possible new TeV-scale particles decaying in the “semileptonic” $\ell\nu qq$ channel, where the resonance decays to a W boson and another W/Z boson, which decay leptonically and hadronically respectively. These results are published in Ref. [8], with no significant excess with respect to the background-only prediction found. Limits are placed on several benchmark new physics models using these results. The search is then widened in Chapter 6, which details the combination of all ATLAS diboson and dilepton resonance search results with 36 fb^{-1} in Ref. [9], further improving on the exclusion limits. To probe models which predict non-resonant new physics, a measurement of the electroweak vector-boson scattering cross-section in the semileptonic channels is described in Chapter 7, which was published in Ref. [10]. Lastly, Chapter 8 details the prospects of diboson resonance searches and vector-boson scattering measurements for the planned upgrade of the LHC, the High-Luminosity LHC, which is projected to start taking 14 TeV data at even higher luminosities in 2026. The prospect studies were published in Ref. [11].

Chapter 2

Theory

*The Standard Model is so complex it would be hard to put it on a T-shirt —
though not impossible; you'd just have to write kind of small.*

—Steven Weinberg

The mathematical formulation of our current understanding of physics at the smallest scale is called Quantum Field Theory (QFT) [12–14]. QFTs provide a simultaneous treatment of both quantum mechanics and special relativity. The QFT which currently provides the best description of the observed interactions of matter is known as the Standard Model (SM). While the SM has been very successful at describing various phenomena, several pieces of evidence indicate it may be a subset of a larger theory. Theories which would supersede the SM are typically denoted as Beyond the Standard Model (BSM).

The aim of this work is to experimentally search for signatures of possible BSM theories in the possible decay mode to W/Z boson pairs. To motivate these searches, this chapter will discuss the details of the SM, as well as several BSM theories which can be probed in this final state. Section 2.1 will provide a terse undergraduate level physics motivation for the study of QFT and Section 2.2 will describe the phenomenology of the SM. Lastly Section 2.3 will motivate various BSM theories and their observable implications. Experts familiar with the field can skip to Section 2.3 without loss.

2.1 Introduction to Quantum Field Theories

As noted earlier, QFT is a relativistic formalism of quantum mechanics. An important distinction with respect to non-relativistic quantum mechanics is that QFT is inherently a study of multi-particle states. A sketch of this can be seen when naively considering the Heisenberg uncertainty principle $\Delta E \Delta t \geq \frac{\hbar}{2}$ and the famous relativistic relation $E = mc^2$. The implication of these two equations together is that in a sufficiently small time window it is possible for the energy of any state, even the vacuum, to fluctuate enough to produce new particles. Thus even the treatment of the propagation of an interacting state through time requires the particle

content and multiplicity to vary during its evolution.

In relativistic theories, space and time are considered on an equal basis as 4-coordinates x^μ with¹ $\mu = 0, 1, 2, 3$, which define a 4 dimensional space-time manifold with metric $g_{\mu\nu}$. Many useful observables in relativity are required to be Lorentz invariant, which means they are invariant under Lorentz transformations Λ . The Lorentz transformations are defined as the group that leave the metric invariant $\Lambda^T g \Lambda = g$. The Poincare group is the full symmetry of space-time including translation invariance.

In non-relativistic quantum mechanics a particle is defined by a state $|\psi(t)\rangle$ and its time evolution is given by the Schrödinger equation $i\hbar \frac{d}{dt} |\psi(t)\rangle = H |\psi(t)\rangle$ where H is the Hamiltonian. The position of the state is given by an operator X . This formulation of quantum-mechanics via the Schrödinger equation is not Lorentz invariant as it treats space and time separately, without restrictions of relativity.

Alternatively in a QFT, the main mathematical objects are fields, which are functions of both space and time, in tensor representations of the Poincare group (e.g. $\phi(\mathbf{x})$, $A_\mu(\mathbf{x})$, $h_{\mu\nu}(\mathbf{x})$ for rank 0, 1, 2 tensors²). In the canonical quantization formulation, a classical field theory is promoted to a QFT by imposing commutation relations on these fields, similar to $[x, p] = i\hbar$ in non-relativistic quantum mechanics. In this perspective the fields are promoted to operators, which are a superposition of ladder operators which can excite or de-excite the particle content of a state. A particle ϕ with definite momentum $\mathbf{p}$ can be made by operating the respective field against the vacuum state $|\Omega\rangle$ (e.g. $|\phi, \mathbf{p}\rangle = \phi(\mathbf{p}) |\Omega\rangle$). The definition of a particle in this context is a set of states $|\psi\rangle$ which transform under some representation of the Poincare group. To represent a relativistic quantum theory, we require the inner product of such states to be Lorentz invariant and unitary. All the unitary representations of the Poincare group were identified by Wigner [15], and can be represented by two numbers, their mass m and spin J which can take any half-integer value ($J = \frac{1}{2}, 1, \frac{3}{2}, \dots$). Particles with fractional spin values are called fermions and those with integer values bosons.

The main model building tool for a QFT is the Lagrangian density $\mathcal{L}$, conventionally just called a Lagrangian, which specifies the particle content of a model and the interactions. Equations of motion can be derived from the principle of least action, which states that $S = \int \mathcal{L} d^4x$ is stationary along a classical path. This results in the Euler-Lagrange Equations, which for a scalar field ϕ are

$$\frac{\partial \mathcal{L}}{\partial \phi} = \partial_\mu \frac{\partial \mathcal{L}}{\partial (\partial_\mu \phi)} \quad (2.1)$$

Similar equations can be derived for any tensor field. Lagrangians can typically be factorized into kinetic terms $\mathcal{L}_{kin}$ which specify how free fields would propagate, and terms $\mathcal{L}_{int}$ which specify the interactions of the fields with each other.

¹We will consider Latin indices $i = 1, 2, 3, \dots$ to index a variety of categories, while we will reserve Greek letters $\mu = 0, 1, 2, 3$ to indicate space-time indices. Lower space-time indices indicate contra-variant components and upper indices as covariant components. Bold letter, such as $\mathbf{x}$, represent 4-vectors in a coordinate free notation. The derivative with respect to space-time is given by ∂_μ .

²We will often suppress the arguments of fields (e.g. $\phi, A_\mu, h_{\mu\nu}$)

Beyond the global Lorentz invariance, Lagrangians can also contain additional internal symmetries. For example, the Lagrangian of a complex scalar field

$$\mathcal{L} = \partial_\mu \phi^* \partial^\mu \phi - m \phi \phi^* \quad (2.2)$$

is symmetric under the transformation $\phi \rightarrow e^{-i\alpha} \phi$. If such a transform is spacetime independent it is called a global symmetry. If it is a function of space-time, $\alpha(x)$, it is then called a local or gauge symmetry. By Noether's theorem, continuous symmetries of the Lagrangian result in conserved currents when the equations of motion are satisfied [16]. Symmetries are often discussed in the context of group theory, where they are identified with a group (e.g Z_2 , $SU(3)$, etc).

An important subset of QFTs are those which have a gauge symmetry, known as gauge theories. For fields with kinetic terms, the presence of an associated gauge symmetry requires that there must be massless spin-1 fields A_μ in the theory associated to the symmetry³ [13]. The spin-1 gauge fields acts as a connection allowing comparison of nearby field points in a gauge invariant way. A local symmetry can be converted to a gauge symmetry by replacing the conventional derivative for the original fields with the gauge-covariant one $\partial_\mu \rightarrow D_\mu = \partial_\mu - ig A_\mu$. The Lagrangian for a single massless spin-1 particle is

$$\mathcal{L} = -\frac{1}{4} F^{\mu\nu} F_{\mu\nu} = -\frac{1}{4} (\partial_\mu A_\nu - \partial_\nu A_\mu) (\partial^\mu A^\nu - \partial^\nu A^\mu) \quad (2.3)$$

which changes under a gauge transform as $A_\mu \rightarrow A_\mu + \partial_\mu \alpha(x)$. Since this change introduces interactions in a an originally free theory, it is conventionally said that the A_μ fields represent gauge bosons which mediate a new force. By Noether's theorem all fields have quantum numbers, sometimes called charges, associated to the induced conserved currents. Matter coupled to a spin-1 particle implies the existence of anti-particles with equal mass and spin but opposite quantum numbers. Gauge theories will form the underpinning of our understanding of the Standard Model in Section 2.2.

One of the fundamental calculations in QFT is the scattering of some initial state particles $|i\rangle$ into some final state particles $|f\rangle$. The differential probability of scattering is given by a formulation of Fermi's Golden Rule

$$dP = \frac{|\langle f| S |i\rangle|^2}{\langle f|f\rangle \langle i|i\rangle} d\Pi \quad (2.4)$$

where S is the scattering matrix which encodes the physics and $d\Pi = \prod_f \int \frac{d^3 p_f}{(2\pi)^3} \frac{1}{2E_f} 2\pi^4 \delta(\sum p_i - \sum p_f)$ is the Lorentz invariant phase-space of the states. Here $\prod_f$ is the product over all final states, p_f/E_f represent the final state momenta/energy, and $\sum p_i$ the total initial state

³The reverse logic can also be held true. Any consistent theory with a massless spin-1 fields requires all fields it interacts with to obey the gauge invariance.

momenta. The S -matrix can be divided into

$$\langle f | S | i \rangle = \langle f | I | i \rangle + (2\pi)^4 \delta(\sum p) \langle f | \mathcal{M} | i \rangle \quad (2.5)$$

where I is the identity matrix relevant for the trivial case of no scattering and the matrix element $\mathcal{M}$ represent the non-trivial scattering of states. Convenient tools known as Feynman diagrams pictorial assist in the calculation of matrix elements from a given Lagrangian through perturbation theory.

An important application of Equations 2.4 - 2.5 is the calculation of the scattering probability, known as the cross-section σ . For the situation with two incoming particles, X_1, X_2 with corresponding energies E_1, E_2 scattering into some new state Y , the cross-section can be calculated as:

$$\sigma(X_1 X_2 \rightarrow Y) = \int \frac{|\mathcal{M}_{X_1 X_2 \rightarrow Y}|^2}{2E_1 E_2 v} d\Pi \quad (2.6)$$

where $\mathcal{M}_{X_1 X_2 \rightarrow Y}$ is the corresponding matrix element for the process, and v is the relative velocity between the two particles (in units of c). Similarly, the decay rate Γ of a particle X to some combination of particles Y can be viewed as a scattering process and can be calculated as

$$\Gamma(X \rightarrow Y) = \frac{|\mathcal{M}_{X \rightarrow Y}|^2}{2m_X} d\Pi \quad (2.7)$$

where $\mathcal{M}_{X \rightarrow Y}$ is again the matrix elements for this process and m_X is the particles mass. The sum of decay rates is known as the width and the relative rate of a specific decay mode with respect to the width is called the branching ratio.

2.2 Standard Model

The Standard Model (SM) is the QFT which is currently our best description of the observed universe. It is a gauge-theory with a local $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ symmetry where the respective gauge groups describe the non-gravitational forces observed in nature. The $SU(3)_c$ group describes the strong nuclear force, which at low-scales binds quarks into hadrons [17]. It is described by the theory of Quantum Chromo-Dynamics (QCD), where the $SU(3)_c$ charge is called colour [18]. For particles in the fundamental representation of this group the colour charge is denoted by one of three abstract values: $c = r, g, b$. The $SU(2)_L \otimes U(1)_Y$ conjointly describe the theory of the weak nuclear force and electromagnetism [19–21]. The conserved quantum numbers associated to the fields are the third component of weak isospin T_3 and the weak hypercharge Y . Due to the Higgs mechanism, the $SU(2)_L \otimes U(1)_Y$ is spontaneously broken to a $U(1)_Q$ gauge symmetry by the vacuum [22–26]. This new symmetry describes Quantum Electro-Dynamics (QED) and the associated quantum number is the electric charge, a combination of weak isospin and weak hypercharge $Q = T_3 + \frac{1}{2}Y$.

QFTs are plagued by multiple theoretical infinities, which require techniques known as

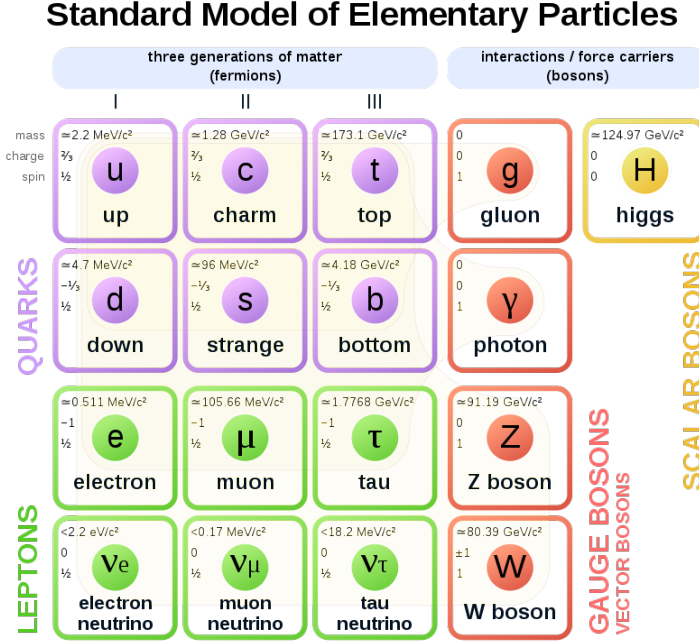


Figure 2.1: The matter content of the Standard Model [30]

renormalization to predict finite experimental results. Theories with which this can be done with only a finite number of model parameters are known as renormalizable and imply a finite number of measurements can be used to define the theory completely. The gauge theories of the SM have been shown to be renormalizable [27]. One observable impact due to renormalization is that the strength of the forces varies with respect to the energy/length scale of the states being probed. For the matter content of the SM it is found that the QCD coupling decreases at high energy, which is known as asymptotic freedom [28, 29]. At high energies the QCD coupling is sufficiently small such that perturbation theory is reliable. At low energies, typically below $O(1 \text{ GeV})$, the coupling constant is large enough that perturbative calculations can not be made. Around this scale a different phenomenon known as colour confinement is observed, where colour charged particles combine to form colour neutral objects. Color neutral bound states of three quarks or anti-quarks are known as baryons, and bound states of a single quarks and single anti-quark are known as mesons. Together, baryons and mesons are known as hadrons.

The particle content of the SM can be seen in Figure 2.1. It can be divided up into two families of spin- $\frac{1}{2}$ fermions, called quarks and leptons, and the gauge bosons associated to the gauge symmetries. Each of the fermion families are further divided into three generations, and by electric charge. Particles in the same family with the same charge but in different generations have the same quantum numbers but differing masses⁴. An additional scalar boson field induced by the Higgs mechanism, known as the Higgs boson, is also included.

The description, phenomenology, and implications of the SM will be developed further in

⁴With the exception of the neutrino masses, which are currently experimentally unknown.

Femion Family	Electric Charge (Q)	Weak Isospin (T_3)	Weak Hypercharge (Y)	Color Representation
u_L, c_L, t_L	$+\frac{2}{3}$	$+\frac{1}{2}$	$+\frac{1}{3}$	3
d_L, s_L, b_L	$-\frac{1}{3}$	$-\frac{1}{2}$	$+\frac{1}{3}$	3
e_L^-, μ_L^-, τ_L^-	-1	$-\frac{1}{2}$	-1	1
ν_e, ν_μ, ν_τ	0	$+\frac{1}{2}$	-1	1
u_R, c_R, t_R	$+\frac{2}{3}$	0	$+\frac{4}{3}$	3
d_R, s_R, b_R	$-\frac{1}{3}$	0	$-\frac{2}{3}$	3
e_R^-, μ_R^-, τ_R^-	-1	0	-2	1

Table 2.1: The fermionic matter content of the SM and the associated quantum numbers with respect to underlying gauge groups. Left and right-handed fermion families are shown separately. The last column indicates the size of the $SU(3)_c$ representation it transforms under (anti-particles of **3** transform under the conjugate representation $\bar{\mathbf{3}}$).

the following sections.

2.2.1 Fermionic Matter Content

Most of the particle content of the SM is composed of spin- $\frac{1}{2}$ fermions. They can be divided up into six quarks (q) which have non-zero quantum numbers under all the gauge groups, and six leptons (ℓ) which have zero colour-charge. The quarks are in the fundamental representation of the $SU(3)_c$ force, and carry one of three colour charges. The six quarks all have distinct masses and are denoted up, down, strange, charm, bottom, top (u, d, s, c, b, t), in order of increasing mass. They can be divided into up-type quarks (u, c, t) with electric charge $Q = \frac{2}{3}$ and down-type quarks (d, s, b) with electric charge $Q = -\frac{1}{3}$. The leptons can similarly be divided into electrically charged leptons, known as electrons, muons, and taus (e, μ, τ), as well as $Q = 0$ leptons known as neutrinos (ν_e, ν_μ, ν_τ). In the SM the neutrinos are all massless. The quantum numbers of all SM fermions are shown in Table 2.1.

A fermion of definite chirality can be represented by a single Weyl spinor $\psi_{L/R}$. Introducing a mass term couples the two chiralities. For massive fermions it is then useful to instead consider bispinors⁵ $\psi = \begin{pmatrix} \psi_L \\ \psi_R \end{pmatrix}$ and the corresponding anti-bispinor $\bar{\psi} = \begin{pmatrix} \psi_R^\dagger & \psi_L^\dagger \end{pmatrix}$.

The Lagrangian for a massive fermion is given by the Dirac Lagrangian [31]

$$\mathcal{L} = i\bar{\psi}(\not{\partial} - m)\psi \quad (2.8)$$

where the slashed notation $\not{\partial}$ denotes the contraction of the 4-vector a_μ with the gamma matrices γ_μ . Due to the chiral structure of the SM, it will also be useful to consider the projection operators $\frac{1-\gamma_5}{2}\psi = \psi_L$ and $\frac{1+\gamma_5}{2}\psi = \psi_R$, where γ_5 is defined as $\gamma_5 = i\gamma_1\gamma_2\gamma_3\gamma_4$.

⁵Spinor indices are almost always suppressed.

Boson	Electric Charge (Q)	Weak Isospin (T_3)	Weak Hypercharge (Y)	Color Representation
$W^\pm$	± 1	± 1	0	1
Z/γ	0	0	0	1
g	0	0	0	8
H	0	$-\frac{1}{2}$	+1	1

Table 2.2: The bosonic matter content of the SM and the associated quantum numbers with respect to underlying gauge groups. The last column indicates the size of the $SU(3)_c$ representation it transforms under.

2.2.2 Gauge Fields

A gauge theory describes a number of massless spin-1 gauge bosons equal to the dimension of the Lie Algebra associated to the group. There are therefore eight $SU(3)_C$ bosons, known as gluons (g), three $SU(2)_L$ bosons (W^i) and a single $U(1)$ boson (B). After spontaneous symmetry breaking discussed in the next section, we will find that the W^i and B fields combine to form the photon γ , and three massive spin-1 bosons $W^\pm$ and Z_0 . The quantum numbers of all SM bosons are shown in Table 2.2.

The Lagrangian for a set of $SU(N)$ gauge bosons is given by the Yang-Mills Lagrangian [32]

$$\mathcal{L} = -\frac{1}{4}F^{a\mu\nu}F_{\mu\nu}^a \quad (2.9)$$

where the field strength tensor $F_{\mu\nu}^a$ is defined as

$$F_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a + gf^{abc}A_\mu^b A_\nu^c \quad (2.10)$$

with the Latin indices⁶ referring to a single component of the Lie Algebra, spanned by the generators T^a . The structure constants, f^{abc} , are group invariants defined by the commutator of the generators, $f^{abc}T^a = [T^b, T^c]$. Lie Algebras with vanishing structure constants, such as $U(1)$, are known as Abelian groups. The g indicates the coupling strength of the fields.

The gauge fields transform under the adjoint representation for an arbitrary gauge transform parameterized by $\alpha^a(x)$

$$A_\mu^a \rightarrow A_\mu^a + \frac{1}{g}\partial_\mu \alpha^a(x) - f^{abc}\alpha^b(x)A_\mu^c \quad (2.11)$$

Fermion fields transform under the fundamental representation of the group they are charged under

$$\psi \rightarrow \exp(-iq\alpha^a(x)T^a)\psi \quad (2.12)$$

where q is the charge of the particle under this group. Explicitly, the generators of the $SU(2)_L$

⁶Occasionally we may suppress the Lie Algebra index which denotes a contraction with the respective generator, $A_\mu = A_\mu^a T^a$.

and the $SU(3)_c$ subgroup in the fundamental representation are the Pauli matrices σ_i , and the Gellman matrices λ_i .

Gauge invariance strictly constrains the form of interactions between fermions and gauge bosons. A gauge invariant alternative of Equation 2.8 can be found by substituting the conventional derivative with a gauge-covariant derivative

$$\partial_\mu \rightarrow D_\mu = \partial_\mu + igA_\mu^a T^a \quad (2.13)$$

This induces linear interactions between fermion and bosons of the form

$$\mathcal{L}_{int} = ig\bar{\psi}\gamma^\mu A_\mu\psi \quad (2.14)$$

From Equation 2.9 and Equation 2.10 it can be seen that non-Abelian groups have both cubic and quadratic couplings between the bosons of the form

$$\mathcal{L}_{A^3} = \frac{g}{2}f^{abc}(\partial_\mu A_\nu^a - \partial_\nu A_\mu^a)A^{b\mu}A^{c,\nu} \quad (2.15)$$

$$\mathcal{L}_{A^4} = \frac{g^2}{4}f^{abc}f^{ade}A_\mu^b A_\nu^c A^{d\mu}A^{e\nu} \quad (2.16)$$

For an Abelian group, the gauge bosons do not couple with each other.

2.2.3 Spontaneous Symmetry Breaking

The mass terms for the fermions and gauge bosons explicitly break the $SU(2)_L \otimes U(1)_Y$ gauge symmetry, so naively one would expect all matter in the SM to be massless. Experimentally though, both fermions and several gauge bosons have a measured mass. Therefore the $SU(2)_L \otimes U(1)_Y$ gauge symmetry must be broken. The Higgs mechanism is a solution to this problem relying on spontaneous symmetry breaking, where the underlying equations are gauge-invariant, but the ground state of the system is not [22–26]. In spontaneous symmetry breaking, the vacuum state which we expand around in perturbation theory does not respect the global symmetry. An expansion around this minima then appears to have a broken symmetry.

The Higgs mechanism introduces a new complex scalar doublet $SU(2)_L$ doublet with $Y = \frac{1}{2}$ denoted ϕ , which can be expanded as:

$$\phi = \begin{pmatrix} \phi^+ \\ \phi_0 \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} \phi_1 + i\phi_2 \\ \phi_3 + i\phi_4 \end{pmatrix} \quad (2.17)$$

The Higgs field is then given self-interactions of the form

$$\mathcal{L} = |\partial_\mu \phi|^2 - \mu^2 |\phi|^2 - \lambda |\phi|^4 \quad (2.18)$$

where μ^2 appears as a wrong-sign mass term and λ denotes some quartic self-coupling. In the

case of $\mu^2 < 0$ and $\lambda > 0$ the potential of ϕ has a minima for:

$$|\phi| = \frac{v}{\sqrt{2}} = \sqrt{-\frac{\mu^2}{2\lambda}} \quad (2.19)$$

where v is the called vacuum expectation value and corresponds to the matrix element of the field in the vacuum $\langle 0|\phi|0\rangle$.

A particular vacuum state breaks the $SU(2)_L \otimes U(1)_Y$ symmetry. Expanding Equation 2.18 around the minima, which we can arbitrary choose to be along ϕ_3 , with $\phi_3(x) = v + \eta(x)$ one would find that the remaining ϕ_i bosons have become massless and η has picked up a mass term. The ϕ_1 , ϕ_2 , and ϕ_4 are the massless Nambu-Goldstone bosons predicted in any theory with a spontaneously broken symmetry [33–35].

Now if the scalar fields are to obey a local gauge symmetry, the derivative must be replaced with gauge covariant ones as in Equation 2.13. This introduces couplings between the scalar bosons and the gauge bosons. We can also choose a gauge, known as the unitary gauge, such that scalar doublet is simplified to

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

where h is the Higgs boson. Utilizing the unitary gauge, the Lagrangian after electroweak spontaneous symmetry breaking is

$$\begin{aligned} \mathcal{L} = & \frac{1}{2} \partial_\mu h \partial^\mu h + \mu^2 h^2 - \lambda v h^3 - \frac{\lambda}{4} h^4 - \frac{1}{4} W_{\mu\nu}^i W^{i,\mu\nu} - \frac{1}{4} B_{\mu\nu} B^{i,\mu\nu} \\ & + \frac{g_{T_3}(v+h)^2}{8} (W_\mu^1 + iW_\mu^2) (W^{1,\mu} - iW^{2,\mu}) + \frac{(v+h)^2}{8} (g_{T_3}W_\mu^3 - g_Y B_\mu) (g_{T_3}W^{3,\mu} - g_Y B^\mu) \end{aligned} \quad (2.21)$$

The bosons can be rewritten into mass eigenstates as

$$W_\mu^\pm = \frac{1}{\sqrt{2}} (W_\mu^1 \pm iW_\mu^2) \quad Z_\mu = \frac{g_{T_3}W_\mu^3 - g_Y B_\mu}{\sqrt{g_{T_3}^2 + g_Y^2}} \quad A_\mu = \frac{g_{T_3}W_\mu^3 + g_Y B_\mu}{\sqrt{g_{T_3}^2 + g_Y^2}} \quad (2.22)$$

The $W_\mu^\pm$ fields have a mass given by $m_W = \frac{g_{T_3}v}{2}$, and Z_μ a mass $m_Z = \frac{g_{T_3}v}{2\cos\theta_W}$, where $\theta_W = \tan^{-1}(\frac{g_Y}{g_{T_3}})$. The A_μ field, now identified with the photon, remains massless and appears as a leftover $U(1)_Q$ gauge symmetry with coupling strength $g_Q = g_{T_3} \sin\theta_W$. The Higgs field also is found to have a positive mass-term $m_h = -2\mu^2$.

The fermions can be given mass terms in a gauge invariant way using the Higgs boson as well. Left-handed fermions transform as $SU(2)_L$ doublets and right-handed fermions as singlets since they have no charge under $SU(2)_L$. We will denote them as

$$L_i = \begin{pmatrix} \nu_{iL} \\ \ell_{iL} \end{pmatrix}, \quad \begin{pmatrix} u_{iL} \\ d_{iL} \end{pmatrix}, \quad R_i = \nu_{iR}, \ell_{iR}, u_{iR}, d_{iR} \quad (2.23)$$

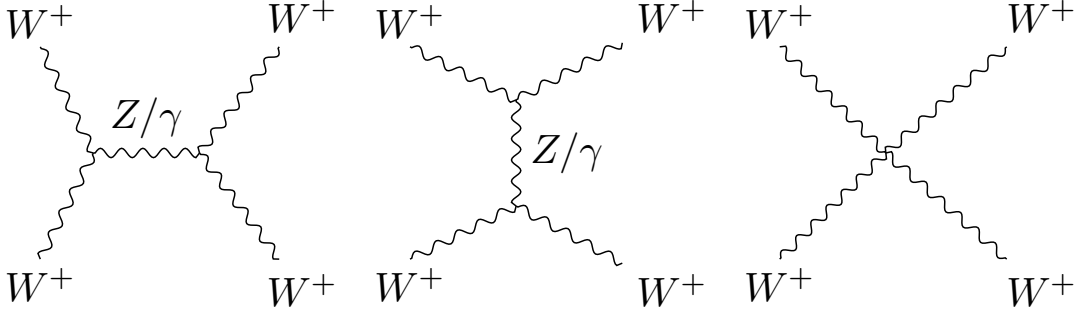


Figure 2.2: Diagrams contributing to same-sign W boson scattering which do not contain the Higgs boson

where i to index the lepton, up-type quark, and down-type quark generations. Both up-type and down-type fermions can be given gauge invariant mass terms by Yukawa couplings with the Higgs bosons

$$\mathcal{L}_{Yukawa} = -y_{ij}^d \bar{L}_i \phi R_j - y_{ij}^u \bar{L}_i i \sigma_2 \phi^* R_j + h.c. \quad (2.24)$$

where $h.c.$ refers to additional terms hermitian conjugate to the previous expression. After spontaneous symmetry breaking and diagonalizing the Yukawa couplings, $y_{ij}^u/y_{ij}^d \rightarrow y_f$, the relabeled fermion fields are mass eigenstates with a mass terms $m_f = y_f v^2/\sqrt{2}$ and couplings with the Higgs field given by $m_f v^2 h \bar{L}_i R_j \sqrt{2}$. These mass eigenstates are the fields discussed in Section 2.2.1.

The Lagrangian for the electroweak currents after spontaneous symmetry breaking in a compact notation suppressing fermion labels is

$$\begin{aligned} \mathcal{L} = & -e A_\mu Q \psi \gamma^\mu \psi \\ & - \frac{e}{\sin \theta_W \cos \theta_W} Z_\mu \left(T_3 \psi \gamma^\mu \frac{1 - \gamma_5}{2} \psi - \sin^2 \theta_W Q \psi \gamma^\mu \psi \right) \\ & - \frac{e}{\sqrt{2} \sin \theta_W} (W_\mu^+ \bar{u}^i \gamma^\mu V_{ij} d^j + W_\mu^+ \bar{\nu} \gamma^\mu \ell) + h.c \end{aligned} \quad (2.25)$$

where the first/second/third lines describe the interaction of the photon/ $Z/W^\pm$ with fermionic matter. The matrix V_{ij} is known as the CKM matrix and encodes the difference between the flavor eigenstates and mass eigenstates.

2.2.4 Vector Boson Scattering

An important phenomenological aspect of the Higgs mechanism can be seen in the scattering of the longitudinal component of electroweak vector bosons (VBS). In a theory without the Higgs boson the relevant Feynman diagrams for same-sign $W^\pm$ scattering are shown in Figure 2.2. Calculating the matrix element without the Higgs boson contribution, one would find that approximately at a center of mass of 1 TeV, the matrix element would be $M \approx 1$. This would

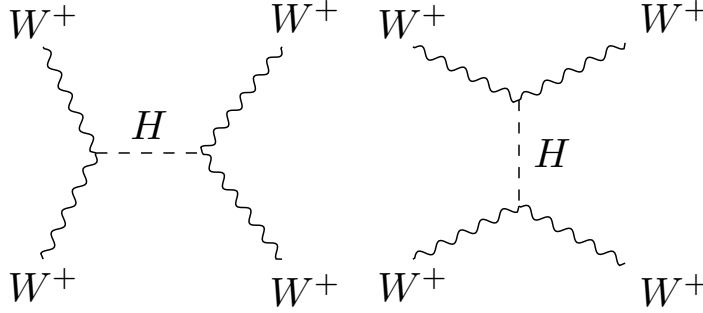


Figure 2.3: Diagrams contributing to same-sign W boson scattering which contain the Higgs boson

imply a violation of unitarity and an indication that perturbation theory in such a theory would break-down.

Including the SM Higgs restores unitarity. The Feynman diagrams for VBS which include the Higgs can be seen in Figure 2.3. This is an aspect of the Goldstone boson equivalence theorem which relates the longitudinal components of the W/Z bosons to the Goldstone bosons produced during spontaneous symmetry breaking. Before its discovery, this process imposed a requirement on the Higgs boson mass known as the Lee-Quigg-Thacker Bound [36]

$$m_h \leq \sqrt{\frac{16\pi}{3}} \frac{1}{v} \approx 1 \text{ TeV} \quad (2.26)$$

If either the Higgs mechanism was not actualized in nature or if the Higgs failed this bound some new non-perturbative physics should have been observable at approximately 1 TeV in this process. These considerations went into the design of the LHC and ATLAS experiment detailed in Chapter 3. The Higgs boson has since been found at a mass of 125 GeV satisfying the unitarity requirements. As indicated by the above discussion, the VBS process is very sensitive to the gauge and Higgs structure of the SM. Any deviations from the SM can induce large changes to the cross-section of this process. A measurement of the VBS process is conducted in Chapter 7.

2.3 Beyond the Standard Model

While the SM has managed to match experimental predictions very accurately, it has both experimental and theoretical shortcomings. Experimental observations which the SM has not been able to describe include the evidence of Dark Matter [37, 38], and the observed matter-antimatter asymmetry [39]. Theoretical issues with the SM include the hierarchy problem, and the inclusion of gravity in a renormalizable way.

Explanation of these phenomena requires physics Beyond the Standard Model (BSM). These theories typically introduce new particles which alter the electroweak sector or introduce ad-

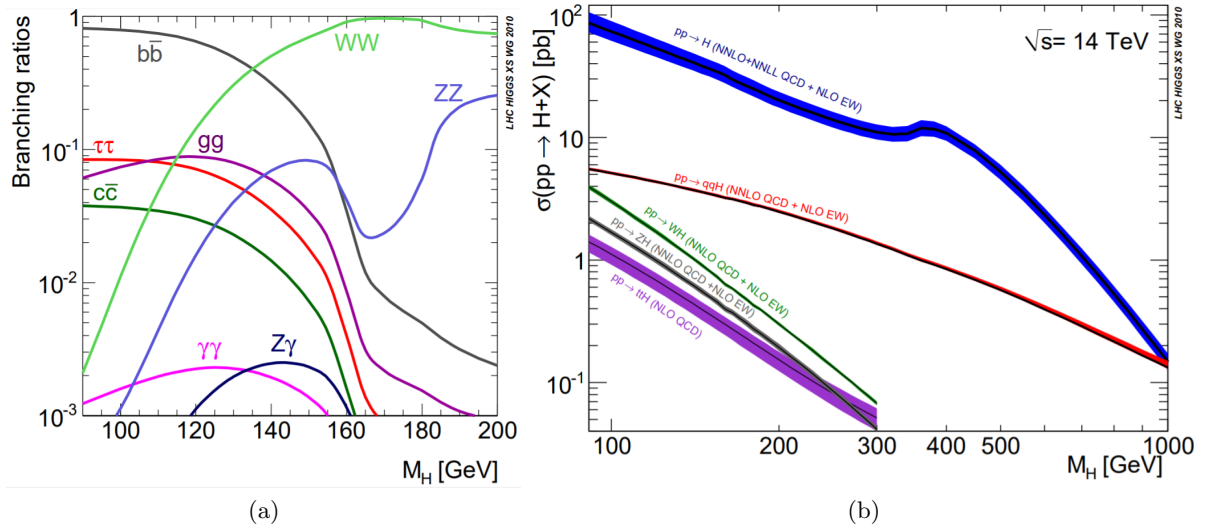


Figure 2.4: The SM Higgs a) branching ratio and b) cross-section as a function of Higgs mass [40].

ditional Higgs Bosons. Many of these new particles are predicted to couple to pairs of W/Z bosons and can decay to these states if they are of sufficient mass. Of specific interest in this work are new particles with masses on the TeV scale. This mass range is light enough to possibly not completely decouple from SM Higgs/electroweak physics but still massive enough to not be constrained by existing measurements. The main production modes of new particles in pp colliders, which are dependent on the spin of the resonance and its model parameters, are typically gluon-gluon fusion (ggF), quark-quark annihilation ($q\bar{q}$), or vector-boson fusion (VBF).

Several BSM theories which attempt to explain various phenomena are described in Sections 2.3.1 - 2.3.2. Direct experimental constraints are placed on a subset of these models in Chapters 5 - 8. Other models may share similar phenomenological features to the benchmark models and the experimental results can be re-interpreted to provide limits for these situations as well.

2.3.1 Extended Higgs Sectors

Several BSM theories predict extensions to the Higgs mechanism which include a spectrum of new scalar particles in addition to the sole one predicted by the SM. Many of these models directly attempt to solve the hierarchy problem. Since the Higgs mechanism is tightly related to the longitudinal components of W/Z bosons, the addition of extra particles in the Higgs sector usually produces observable differences in SM Higgs decay modes and W/Z related cross-sections. Frequently the new Higgs Sector particles are predicted to decay on-shell to W/Z bosons.

Several concrete examples of extended Higgs sectors are discussed below. We will probe

such models by approximating the new resonant states as having SM-like Higgs couplings but with heavier masses and widths narrower than detector resolution effects. The branching ratios and cross-sections for this toy model can be taken from SM Higgs predictions as displayed in Figure 2.4.

Singlet Extension

The simplest extension of the Higgs sector involves the addition of a new scalar $SU(2)$ -singlet field S which does not acquire a vacuum expectation value [41, 42]. The most general Lagrangian of such a scalar only interacting with the SM Higgs sector is [41]

$$\begin{aligned}
 -\mathcal{L} = & \frac{m^2}{2} h^\dagger h + \frac{\lambda}{4} (h^\dagger h)^2 + \frac{\delta_1}{2} h^\dagger h S + \frac{\delta_2}{2} h^\dagger h S^2 \\
 & + \frac{\delta_1 m^2}{2\lambda} S + \frac{\kappa^2}{2} S + \frac{\kappa^3}{3} S^3 + \frac{\kappa^4}{4} S^4
 \end{aligned} \tag{2.27}$$

where the δ 's and κ 's are new model parameters. In the most general case, the SM Higgs and S kinematically mix, with the newest lightest state retaining the SM Higgs branching ratios but with reduced couplings proportional to the $\cos^2(\phi)$, where ϕ is the mixing parameter. To be consistent with current Higgs measurements, such a model would need to be near the decoupling limit where $\phi \rightarrow 0$. The decay width of the heavier state to SM particles is suppressed by $\sin^2(\phi)$.

New scalars which do not mix with the SM Higgs are also predicted in some theories. In such models the new scalars have direct coupling operators with the rest of the SM particles and represent some new sector of physics not explicitly connected to the Higgs. In Ref. [43] all possible spin-0 states in arbitrary $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ representations which have possible linear renormalizable interactions with SM particles are classified.

Two-Higgs Doublet Models

The next most complicated model is the addition of a second Higgs $SU(2)$ -doublet, known as a Two-Higgs Doublet Model (2HDM) [44]. While such models initially seem more complicated, they are motivated by supersymmetric theories, which required 2HDM, and various models of baryogenesis, as they provide natural sources of CP violation.

For two doublets of hypercharge $Y = +1$, here denoted Φ_1 and Φ_2 , the most generic Lagrangian with only quadratic terms even in the fields and unbroken CP conservation in the Higgs Sector is

$$\mathcal{L} = m_{11}^2 \Phi_1^\dagger \Phi_1 + m_{22}^2 \Phi_2^\dagger \Phi_2 - m_{12}^2 (\Phi_1^\dagger \Phi_2 - \Phi_2^\dagger \Phi_1) + \frac{\lambda_1}{2} (\Phi_1^\dagger \Phi_1)^2 + \frac{\lambda_2}{2} (\Phi_2^\dagger \Phi_2)^2 \tag{2.28}$$

$$+ \lambda_3 \Phi_1^\dagger \Phi_1 \Phi_2^\dagger \Phi_2 + \lambda_4 \Phi_1^\dagger \Phi_2 \Phi_2^\dagger \Phi_1 + \frac{\lambda_5}{2} \left((\Phi_1^\dagger \Phi_2)^2 + (\Phi_2^\dagger \Phi_1)^2 \right) \tag{2.29}$$

As in the ordinary Higgs mechanism, three of these fields provide the W/Z boson their

mass during spontaneous symmetry breaking. This leaves five Higgs states: two neutral scalars h/H , one pseudo-scalar A , and charged scalars $H^\pm$. The two main phenomenological parameters of the 2HDM are the α mixing angle between the neutral scalars and the β mixing angle of the charged scalars and pseudo-scalar. Along with the diagonalized mass values $(m_h, m_H, m_A, m_{H^\pm})$ this is enough to completely specify a 2HDM model.

The general 2HDM framework produces flavour changing neutral currents at tree-level due to the separate Yukawa couplings for each doublet producing possible mixing between flavour states and mass states. If tree-level flavour changing neutral currents exist, they must be highly suppressed to match current experimental observations. They can be removed in a 2HDM by applying additional discrete symmetries, or equivalently a specification of which doublet each right-handed fermion solely couples to. In a Type-I 2HDM, all right-handed fermions couple to the same doublet, while in Type-II the up-type and down-type right-handed fermions couple to different doublets. A priori there is no need for the leptons to follow the same convention. A separate “lepton-specific” 2HDM has all the right-handed quarks and leptons coupling to separate doublets, and a “flipped” 2HDM which is the same as the Type-II but all the leptons coupling to the same doublet as the up-type quarks.

The different 2HDM models are all phenomenologically distinct. Interestingly, for Type I and II models, a state similar to the SM Higgs can be reproduced as $H_{SM} = h \sin(\alpha - \beta) - H \cos(\alpha - \beta)$. Similar to singlet extensions, for Type I and Type II 2HDM models to match Higgs measurements requires that they are near the limit $\cos(\alpha - \beta) \rightarrow 0$. In general for all 2HDM models, the WW and ZZ couplings of the lightest h are the same as the SM Higgs, but proportional to $\sin(\alpha - \beta)$. Models with $\sin(\alpha - \beta) \rightarrow 1$ can be reproduced when the vacuum expectation value aligns only along one of the neutral Higgs components, known as the alignment limit, or when the mass scale of one scalar is significantly smaller than the others, known as the decoupling limit. Current Higgs measurements constrain all 2HDM models to be near these possible limits, which suggests that the branching ratio of the new heavy scalars to VV to be negligible, but there exists a range of parameter space where this process can still dominate [45]. For this space of 2HDM models the main interest is when the new states are larger than the SM Higgs mass but less than 1 TeV. This discussion focuses on the neutral scalar states as the charged and pseudo-scalar states do not have direct diboson decay modes⁷.

Composite Higgs Models

The main aspect of the hierarchy problem arises from the fact that Higgs mass as described by the SM is very sensitive to radiative effects, and hence should be at the effective cut-off scale of the SM. Cancellations of these radiative effects with renormalization typically involve a so-called “fine-tuning” as the cancellation would need to be done to 28 orders of magnitude to produce the observed Higgs mass. Such a problem is addressed in Composite Higgs Models where the Higgs is no-longer a fundamental particle, but a bound state of some new strongly-

⁷Except Vh type decays which are not of interest in this work.

interacting physics below some scale Λ [46, 47].

The general feature of these models is a prediction of a new sector of fermions charged under a new strongly-interacting gauge group $\mathcal{G}$ in addition to the $SU(2)_L \otimes U(1)_Y$ of the SM. At some scale Λ , the new fermions form bound states which break the overall symmetry to a new gauge group $\mathcal{H}$ containing $SU(2)_L \otimes U(1)_Y$ as a sub-group. The coset $\mathcal{G}/\mathcal{H}$ contains pseudo-Nambu-Goldstone bosons, with four required to be in the $SU(2)_L \otimes U(1)_Y$ representation matching that of the SM Higgs doublet [47]. At the scale of the ordinary Higgs mechanism, these states spontaneously break the $SU(2)_L \otimes U(1)_Y$. A frequently studied “minimal” choice is $\mathcal{G} = SO(5)$, $\mathcal{H} = SO(4)$ [47, 48].

Current Higgs and electroweak precision measurements typically constrain such models to the scale $\Lambda \gtrsim 1$ TeV [49]. Directly observable states of Composite Higgs models would be the next lowest composite state, denoted a spin-1 ρ' , which would decay to W/Z pairs and be on the scale of Λ .

Interestingly, it is found that some Composite Higgs models are dual to Bulk Randall-Sundrum Models discussed in Section 2.3.3 via the AdS/CFT correspondence [48].

2.3.2 Extended Gauge Sectors

It is very common in BSM theories to predict new TeV scale spin-1 particles. These can originate from models with extended gauge sectors, such as GUT theories [51, 52], as Kaluza-Klein (KK) modes of SM W/Z bosons from warped extra dimension, as in Section 2.3.3, or from models with strongly coupled sectors, such as in composite Higgs models discussed in Section 2.3.1.

A convenient phenomenological tool for studying a broad range of spin-1 resonances, is the Heavy Vector Triplet (HVT) model [50] which is based on a Simplified Model Lagrangian approach [53]. The underpinning idea is that resonant searches are not sensitive to most of the parameters of a given model, just those related to the on-shell resonant kinematics. For example, resonance searches which are limited by detector resolution have little to no sensitivity to model parameters which affect the tails of the resonance. To study these models more generically, a phenomenological Lagrangian is derived with an Effective Field Theory expansion assuming some new resonant particle content, keeping only terms relevant to on-shell production. The parameters of this Simplified Model Lagrangian are experimentally measurable and can be matched to operators in a specific theory.

The HVT model adds a set of new $SU(2)$ fields with zero hypercharge, denoted V^a , $a = 1, 2, 3$, which describe new charged $V_\mu^\pm$ and neutral V_μ^0 vector states in the same style as the SM W/Z bosons [50]. We will often denote these as W' and Z' states respectively. The Lagrangian

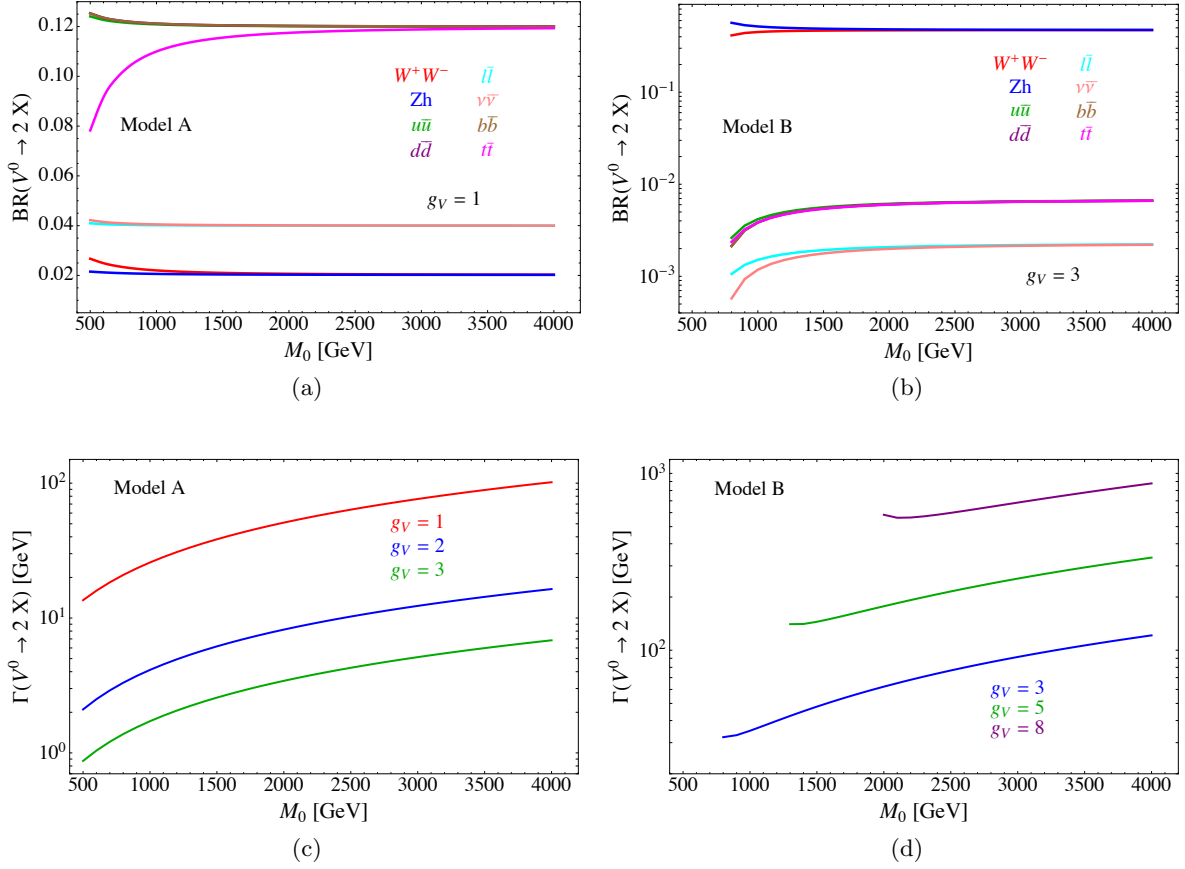


Figure 2.5: The (top) branching ratio and (bottom) total width of the new V^0 boson as a function of mass for the (left) weakly coupled Model A and (right) strongly-coupled Model B parameter choice [50]. For coupling convention see Footnote 8.

of this model is

$$\begin{aligned}
\mathcal{L} = & -\frac{1}{4}D_{[\mu}V_{\nu]}^a D^{[\mu}V^{\nu]}_a + \frac{m_V^2}{2}V_\mu^a V^{\mu a} \\
& + ig_H V_\mu^a \left(H^\dagger \frac{\sigma^a}{2} D_\mu H - D_\mu H^\dagger \frac{\sigma^a}{2} H \right) + g_F V_\mu^a \sum_f \bar{f}_L \gamma^\mu \frac{\sigma^a}{2} f_L \\
& + g_{VVV} \epsilon_{abc} V_\mu^a V_\nu^b D^{[\mu}V^{\nu]}_c + g_{VVHH} V_\mu^a V^{\mu a} H^\dagger H - g_{VVVH} \epsilon_{abc} W^{\mu\nu a} V_\mu^b V_\nu^c
\end{aligned} \tag{2.30}$$

where the new parameters $g_H, g_F, g_{VVV}, g_{VVHH}$, and g_{VVVH} describe the couplings of the new bosons to SM particles⁸. Being explicit, the first line of Equation 2.30 represents the kinematic term of the V fields and its couplings with the W/Z bosons through the covariant derivative

$$D_\mu V_\nu^a = \partial_\mu V_\nu^a + \epsilon^{abc} W_\mu^b V_\nu^c \tag{2.31}$$

The second line contains the direct interaction with SM fermions, as well as with the Higgs boson, which after electroweak symmetry breaking provides decay channels to dibosons through their longitudinal polarization. The coupling to fermions g_F can also be split to different coupling for each fermion. It is also worth noting that new fields will kinematically mix with the SM W/Z bosons after symmetry breaking. The g_{VVV} , g_{VVHH} , and g_{VVVH} terms are not relevant for single V production so will be ignored for the rest of the discussion.

In order to maintain the SM W/Z masses and the relation $\frac{m_W^2}{m_Z^2 \cos(\theta_W)} = 1$ to within experimental precision, it is required that the V masses after symmetry breaking be a factor 100 times greater than $m(W/Z)$ and be nearly degenerate. In this situation the mixing angles are small and Equation 2.30 can be used for calculating observables after symmetry breaking as well.

Several different benchmark coupling values are defined to probe various regimes of possible parameter space. The first of these, denoted Model A, directly corresponds to a model with an extended gauge sector $SU(2)_1 \otimes SU(2)_2 \otimes U(1)_Y$ which gets broken to the SM $SU(2)_L \otimes U(1)_Y$ [54]. The second model, Model B, is a mapping to the linear effects of a composite $SO(5)/SO(4)$ Higgs model [55]. These respectively correspond to weakly and strongly coupled situations of the V boson with the SM bosons. For probing the VBF production modes, a third Model C is used with all fermion couplings set to $g_F = 0$, and $g_H = 1$.

In this formalism the cross-sections of single V production and decay rates can be calculated analytically from the coupling and mass parameters, which are provided in Ref. [50]. The branching ratios and widths of the V^0 boson with the Model A and Model B coupling can be seen in Figure 2.5. Figure 2.6 shows the $q\bar{q}$ production cross-section as a function of g_F and g_V couplings.

⁸A different convention for the couplings constant is used here than Ref. [50], which uses dimensionless constants $c_H, c_F, c_{VVV}, c_{VVHH}, c_{VVVH}$, dimensionful constant G_V and the electroweak coupling constant g . The mapping between conventions follows: $g_H = g_{VCH}$, $g_F = \frac{g_{EWK}}{g_V} c_F$, $c_{VVV} = \frac{g_V}{2} c_{VVV}$, $g_{VVHH} = g_V^2 c_{VVHH}$ and $g_{VVVH} = \frac{g}{2} c_{VVVH}$

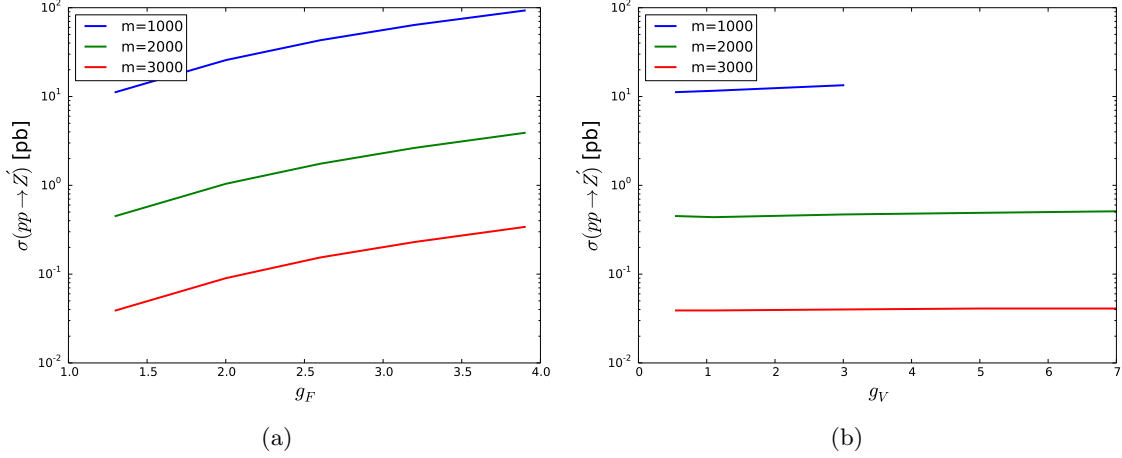


Figure 2.6: The $q\bar{q}$ production cross-section as a function of a) g_F and b) g_H for various Z' masses.

2.3.3 Warped Extra Dimensions

The hierarchy problem can also be solved by the addition of extra spatial dimensions. Gravity would then propagate along all dimensions while the SM interactions are confined to our observable 4-dimensional (4D) spacetime surface, resulting in the disparity of the strength of gravity with respect to the other forces. In such models the extra dimension is typically compact and small in scale so as to be not observable [56, 57]. The 4D Plank-scale is then related to the full $4 + nD$ Plank-scale by the physical scale of the compactified dimension R

$$M_{PL}^2 = M_{PL,4+n}^{2+n} R^n \quad (2.32)$$

Thus the hierarchy problem can be solved if $M_{PL,4+n}$ is near the electroweak scale. There is then no longer a hierarchy of energy but the appearance of one in our 4D universe. These extensions would modify the $1/r^2$ law of gravity below the distance scale R . Since the accuracy of this law is measured to be true at the scale of $O(10) \mu\text{m}$ [58, 59], Equation 2.32 indicates that $n = 1$ additional dimensions are already excluded.

An alternative solution of this form is to not allow the extra dimension to be factorizable and compact. Instead it could be an additional finite dimension with an anti-de Sitter space geometry. Theories of this type are known as Randall-Sundrum (RS) models [60, 61]. The total 5-dimensional metric for RS models is:

$$ds^2 = e^{-2kr_c|\phi|} \eta_{\mu\nu} dx^\mu dx^\nu + r_c^2 d\phi^2 \quad (2.33)$$

where $0 < \phi < \pi$ is the new co-ordinate, πr_c defines the size of the new dimension and k is the warp factor. The hierarchy problem is then addressed by the fact energy scales between the 4-dimensional theory and full 5-dimensional theory are warped by the exponential factor $e^{-\pi k r_c}$.

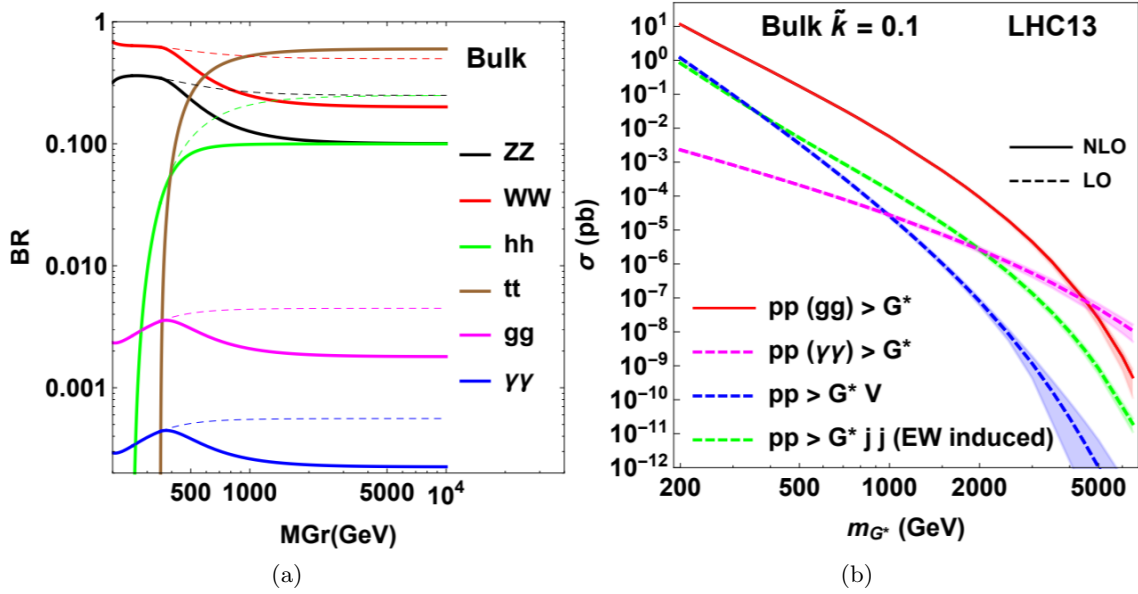


Figure 2.7: The Bulk Bulk RS graviton a) branching ratio and b) cross-section as a function of mass [64].

Values of $\pi k r_c \sim 36$ are consistent with solving the hierarchy problem. The RS equivalent of Equation 2.32 is now

$$M_{Pl}^2 = \frac{M_{PL,5}^3}{k} \left(1 - e^{-2\pi k r_c}\right) \quad (2.34)$$

In the RS framework the metric tensor fluctuations in the usual 4-dimensional space can be viewed as spin-2 gravitons, while the fluctuation in the extra dimension corresponds to a spin-0 field, known as the radion, which is massless in the simplest scenario. A fundamental problem in the original RS framework is that it lacks a mechanism to stabilize the size of the extra dimension, which is also the vacuum expectation value of the radion field. A mechanism which solves both the stability of the warped dimension and the lack of observation of additional massless particles is to introduce an additional bulk scalar, which has its interactions localised on the two ends of the extra dimension [62, 63]. This dynamically produces a radion vacuum expectation value and causes the radion field to acquire a mass term, which is typically smaller than the graviton mass.

In the original formulation of the RS model, the SM fields are constrained to the 4D plane and only the graviton and radion propagate through the bulk. Since the bulk is finite in the 5th dimension, the wave-functions can be decomposed into the 4D fields and the Fourier modes in the additional dimension. The visible particle content would then be an infinite tower of modes of increasing mass for the graviton and radion, called Kaluza-Klein (KK) modes. The mass scale of the first modes would be on the weak-scale.

The original RS model results in predictions which are inconsistent with constraints from flavour physics and electroweak precision measurements. To circumvent this, the SM fields

can be allowed to propagate into the 5th dimension bulk, which suppressed the coupling of the graviton and radion to quarks [65]. This is known as the Bulk RS model. Such models also predict KK modes of all SM particles with higher masses. The graviton decay modes in the Bulk RS model are predominantly to top-quark pairs, di-Higgs, and W/Z bosons [64]. The production cross-section and branching ratio for the graviton can be seen in Figure 2.7.

The radion coupling to fermions is proportional to their mass while it is proportional to the square of the masses for bosonic fields. Phenomenologically the radion is then similar to a Higgs boson, but with a larger mass and smaller width [64]. Results which put constraints on heavy Higgs-like models are also applicable to radions as well. The production cross-section and branching ratio for the radion can also be predicted from Figure 2.4.

The overall phenomenology of the RS models, besides the new KK masses, are dictated by two free parameters: the size of the extra dimension r_c and the warp factor k . When focusing on the graviton phenomenology the only relevant parameters are the resonant mass and the warp factor, often specified in units of the Planck mass k/M_{Pl} [64]. For probing RS models, we will directly search for spin-2 bulk gravitons, while other resonances in this model, such as the radion or the KK gauge bosons, can be reinterpreted from the corresponding spin-0 Higgs-like and spin-1 gauge boson searches.

2.3.4 Effective Field Theories

BSM theories with new high-mass particles can also be probed in a more model independent fashion through use of an effective field theory (EFT) framework. In such a framework the new high-mass physics is integrated out of the Lagrangian, leaving new operators involving only the low-scale particles of the theory. A search for new particles is then related to measurements of the couplings of existing particles.

The SM can be extended to an EFT by the addition of new higher-dimensional operators⁹

$$\mathcal{L}_{EFT} = \mathcal{L}_{SM} + \sum_i \frac{c_i^5}{\Lambda} \mathcal{O}_i^5 + \sum_i \frac{c_i^6}{\Lambda^2} \mathcal{O}_i^6 + \dots \quad (2.35)$$

where Λ the energy scale of the new-physics, and $\mathcal{O}_i^d$ are new operators of dimension d with corresponding Wilson coefficients c_i^d . The EFT is valid as long as the energy scale being studied is below Λ , and the SM is reproduced when $\Lambda \rightarrow \infty$. The operators $\mathcal{O}_i^d$ can include any possible combination of SM fields or Lorentz objects (e.g. derivatives ∂_μ) of mass dimension d . A fundamental aspect of an EFT is that the new operators are power-suppressed by the scale Λ , which is defined to be large in the regime of validity, and hence the new physics amounts to small corrections to the SM at low-energy. For a specific BSM theory, such as those in Sections 2.3.1 - 2.3.3, one can integrate out the new states resulting in a mapping between the original model parameters and the coefficients c_i^d , which can be evolved to the low-scale physics

⁹The SM contains operators of dimension four and a single operator of dimension two, the Higgs mass. Operators of dimension less than five will not be considered part of the EFT expansion.

by the renormalization group equations [66, 67].

It is typically assumed that the new operators obey the $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ gauge symmetry of the SM. Such an EFT is typically known as the SMEFT [66, 67]. There is then only a single $d = 5$ dimensional operator, which results in Majorana masses for the neutrino and lepton number violation [68]. At $d = 6$, 63 independent operators arise, 59 if baryon number conservation is enforced [69]. At this point the expansion provides observable differences to a large number of processes, but usually only a handful of operators contribute to each. In the context of VBS, the operators which contribute would result in alterations to the measured triple-gauge couplings. These operators have been measured by both ATLAS and CMS to be consistent with zero in various searches [70–73]. If these operators are indeed zero, the next contribution to VBS would be with operators which adjust the quadratic gauge coupling, which would arise at $d = 8$. In Ref. [74] the 23 $d = 8$ operators which can alter the VBS process are categorized. The VBS cross-section measurements in Chapter 7 could be used to provide constraints on these operators.

Summary

Our current understanding of nature at the smallest scale is through Quantum Field Theory. The current theory which explains the observed world is known as the Standard Model, with matter content described by several fermion fields and vector bosons fields which mediate the forces. An underlying aspect of the Standard Model is that it is a gauge theory described by a $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ symmetry. The mass terms in the theory can be provided in a gauge invariant way through the Higgs mechanism, which introduces an additional scalar boson into the theory.

The Lagrangian for the SM in a highly condensed notation is:

$$\begin{aligned} \mathcal{L} = & -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} + i\bar{\psi}\not{D}\psi \\ & + D_\mu\phi D^\mu\phi - \frac{m}{2}\phi^2 + \frac{\lambda}{4}\phi^4 \\ & + y_{ij}^d\bar{L}_i\phi R_j + y_{ij}^u\bar{L}_i i\sigma_2 h\phi^* R_j + h.c. \end{aligned} \tag{2.36}$$

where the first line includes the relevant kinetic and interaction terms between the fermions and gauge fields, the middle line includes the Higgs mechanism, and the last line encodes the Yukawa couplings of the Higgs to fermions.

While the predictions of the Standard Model have been able to match experimental observations to high precision, there are several experimental features which it can not explain, such as the existence of Dark Matter and the observed matter-antimatter asymmetry. In addition, there are several theoretical shortcomings of the theory of the Standard Model such as the hierarchy problem which indicate it might not be a complete theory of nature. These facts indicate that there may be some physics Beyond the Standard Model. Such new models are very diverse

in how they address the issues of the Standard Model and have been studied extensively in the literature. A common feature of these new models is the prediction of new physics which alter the dynamics of the production of pairs of electroweak vector bosons. This can originate from new particles which can decay on-shell to these states, or alterations of the Standard Model couplings at some energy scale. Direct experimental constraints on these models from ATLAS data in final states with pairs of SM W/Z bosons are presented in Chapters 5 - 8.

Chapter 3

The ATLAS experiment

It's still magic even if you know how it's done.

—Terry Pratchett, *The Wee Free Men*

The main observables of a QFT as discussed in Chapter 2 are the scattering cross-sections of processes and their kinematics. One of the historically most successful ways to measure these observables is collider experiments, where specific initial states particles are collided at high energies. Due to the nature of a QFT, the out-going particles and their multiplicities can be different from the initial states. Thus, collider experiments are designed to be able to detect and identify particles of different types in order to reconstruct the whole underlying interaction.

The ATLAS experiment is built around a general purpose detector designed to reconstruct and identify particles from the collision of pairs of protons. For this purpose, it is composed of a series of tracking systems and calorimeters, which measure the trajectories and energy of the outgoing particles. The proton collisions analyzed by the ATLAS experiment are produced by the Large Hadron Collider, which accelerates the protons to 6.5 TeV. The following sections summarizes the LHC machine in Section 3.1 and the ATLAS experiment in Section 3.2.

3.1 Large Hadron Collider

The Large Hadron Collider [75–78] (LHC) is a double-ring superconducting particle accelerator that spans 26.7 km in circumference through the French-Swiss countryside operated by The European Organization for Nuclear Research (henceforth denoted CERN). The LHC is designed to accelerate two beams of charged hadrons in opposite directions which necessitates the double-ring structure as the two beams need to be bent by opposite magnetic fields. Due to space constraints in the LHC tunnel, which was inherited from the LEP experiment, the two beam-pipes are stored in the same support system, approximately 194 mm apart, with a twin-bore magnet design applying dipole magnetic fields on the respective beams. The beam

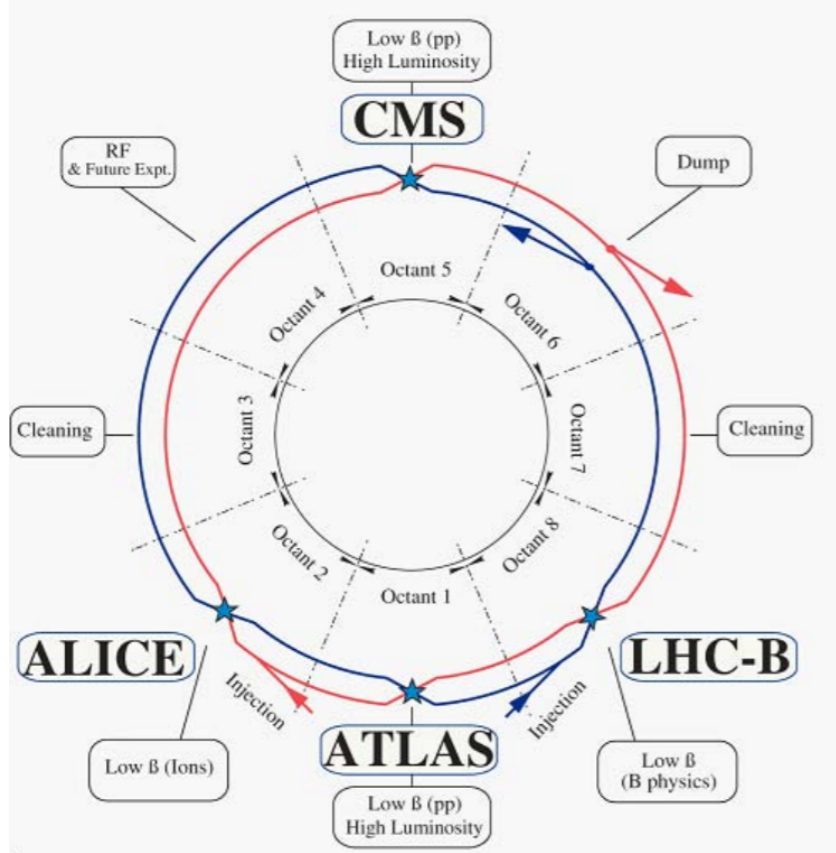


Figure 3.1: Picture of the LHC ring in octants alongside the specialized equipment housed in the respective straight-sections [75].

momentum in GeV for a circular particle accelerator is determined by

$$p = 0.3 B \rho \quad (3.1)$$

where B is the magnetic field strength and ρ is the radius of curvature. The peak magnetic field required for 7 TeV proton (p) beams with the LHC geometry is then 8.33 T, which is provided by NbTi superconducting magnets cooled to 1.9 K. The cooling is done via superfluid helium. To stabilize the beams, a strong focusing technique is used where alternating quadrupoles focusing individually on the horizontal and vertical directions provide stronger focusing overall. The dipole bending magnets are placed intermediately between the quadrupoles. Each LHC lattice cell is 106.9 m long and includes two quadrupoles and six bending dipoles. Higher order sextupole and decapole magnets are also placed in the lattice to correct for chromatic aberrations and other effects.

Superconducting radio frequency (RF) cavities operating at 400 MHz are used to accelerate the beams up to 6.5 TeV. Each individual RF system has eight 36 cm long cavities producing 2 MV of accelerating voltage over their length. A separate klystron is used to power each cavity individually.

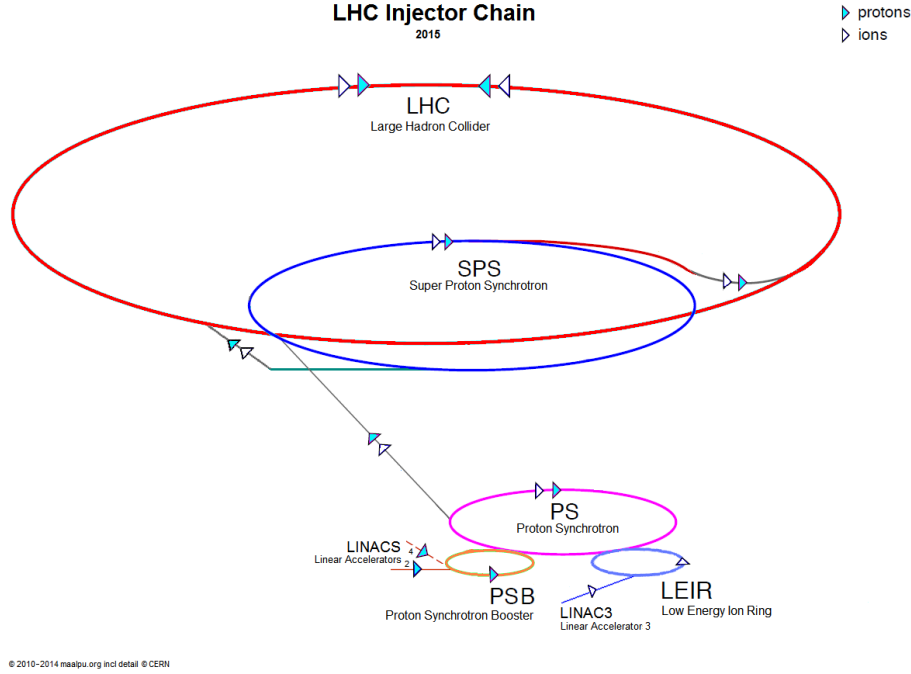


Figure 3.2: The LHC injection chain [83].

The overall geometry of the LHC is divided into octants with an approximately 528 m straight section in the middle. Each of these 8 straight sections correspond to possible interaction regions (IR), but only four of those are instrumented for physics experiments. The design purpose of each of the LHC IRs can be seen in Figure 3.1. The four physics experiments located at the colliding IRs are:

- IR1: The ATLAS general purpose experiment [79]. The ATLAS detector is further detailed in Section 3.2
- IR5: The CMS general purpose experiment [80].
- IR2: The ALICE experiment designed to study heavy-ion collisions [81].
- IR8: The LHCb experiment which is designed to study b -physics [82].

The other four remaining IRs are not designed for colliding beams but instead they house other systems required for operation. IR3 and IR7 contain collimation systems to reduce the spread in the beam momentum and spatial size respectively. The RF acceleration system described above is housed in IR4 for each beam. IR6 holds the kicker magnets which dump the beam out of the LHC in case of emergency or at the end of an operation cycle. In addition to the dedicated experiments, IR2 and IR8 also house the beam injection systems.

To supply the hadrons for LHC operation, a chain of accelerators inherited from the LEP machine is used and displayed in Figure 3.2. For protons the chain follows:

- LINAC2: Ionizes hydrogen gas and accelerates it to 50 MeV.
- Proton Synchrotron Boosted (PSB): Further accelerates these protons to 1.4 GeV. One fill of the LINAC2 is defined as a PSB bunch.

- Proton Synchrotron (PS): Accepts 6 bunches from the PSB, splits them into 72, and accelerates them to 25 GeV. At this point the bunches are separated by 25 ns.
- Super Proton Synchrotron (SPS): Further accelerates the bunches from the PS up to 450 GeV and fills the LHC up to maximum operating value of 2808 bunches.

A proton bunch in the LHC is typically composed of approximately 10^{11} protons.

3.1.1 Luminosity Measurements

For a specific process with a cross-section σ_{proc} , the number of events expected to be produced in a collider experiment is given by

$$N_{proc} = \int dt \mathcal{L} \sigma_{proc} \quad (3.2)$$

where $\mathcal{L}$ is the instantaneous luminosity and the integral is over time. The instantaneous luminosity is related to the beam-parameters and not on the underlying physical process being considered. Since cross-sections are independent of time, one can define the integrated luminosity as $L = \int dt \mathcal{L}$, which simplifies Equation 3.2 to

$$N_{proc} = L \sigma_{proc} \quad (3.3)$$

In order to study processes with very low cross-section it is important for an experiment to produce as high instantaneous/integrated luminosity as possible. This is one of the physics goals of the LHC.

The instantaneous luminosity of a circular proton accelerator can be calculated from the accelerator and collision parameters as

$$\mathcal{L} = \frac{N_b^2 n_b f_{rev} \gamma}{4\pi \epsilon_n \beta^*} F \quad (3.4)$$

where N_b is the number of protons in a bunch, n_b is the number of bunches in the accelerator, f_{rev} is the revolution frequency, γ is the relativistic factor, ϵ_n is the normalized emittance, β^* is the β -function at the interaction point, and F is a geometric reduction factor. The instantaneous luminosity is not constant over an accelerator run and decreases exponentially due to the particle collisions and experimental beam losses. The typical run length of the LHC is approximately 20 hours per fill.

One detriment to a high luminosity machine is that the number of collisions in a bunch crossing, denoted in-time pile-up, is greater than one. These collisions are predominately inelastic pp collisions, which are mostly uninteresting for physics studies. These pile-up events can significantly reduce physics performance such as triggering, object reconstruction, and calibrations. The pile-up parameter μ is related to the bunch-by-bunch instantaneous luminosity

Machine Parameter	Design	2015	2016
Beam Energy [TeV]	7	6.5	6.5
Typical N_b [10^{-11}]	1.15	1.1	1.1
n_b	2808	2232	2208
Bunch Spacing [ns]	25	25	25
f_{rev} [kHz]	11.245		
β^* [cm]	55	80	40
Average $\langle\mu\rangle$	23	13	25
Peak $\mathcal{L}$ [10^{34} cm $^{-2}$ s $^{-1}$]	1	0.5	1.3
L [fb $^{-1}$]	N/A	4.0	38.5

Table 3.1: The LHC beam parameters at design and during 2015 - 2016 operations as measured by the ATLAS experiment [86, 87].

by [84]

$$\mathcal{L}_b = n_b \frac{\mu f_{rev}}{\sigma_{inel}} \quad (3.5)$$

where σ_{inel} is the pp inelastic cross-section, approximately 80 mb. The pile-up parameter is usually averaged over intervals of approximately 1 minute in length, called a lumi-block, and is denoted $\langle\mu\rangle$. The effect of additional collisions from neighboring bunch-crossing can also be noticeable, especially for detector elements with slow responses, and is denoted as out-of-time pile-up.

The BCM and LUCID detectors (described in Section 3.2.4) of the ATLAS experiment are designed to measure the luminosity at the ATLAS collision point, but can only measure the visible fraction of these collisions within their acceptance. When considering only the visible cross-section σ_{vis} , Equation 3.5 becomes

$$\mathcal{L}_b = n_b \frac{\mu_{vis} f_{rev}}{\sigma_{vis}} \quad (3.6)$$

where μ_{vis} is the visible pile-up parameter. Both σ_{vis} and μ_{vis} depend on the exact definition of the detector acceptance, which is usually divided into several sub-detectors. A specific luminosity algorithm is a choice of a subset of these sub-detectors.

To calibrate this procedure, the value of σ_{vis} needs to be derived for each algorithm. This is done with a technique known as a Van der Meer scan [85]. For calibration of the Van der Meer scan method, dedicated runs with high β^* and low amounts of fills are operated [84, 86]. The derived calibrations are transferred to nominal run parameters by comparisons between algorithms with uniform performance over different parameters.

The LHC parameters used for operation during the years 2015 - 2018, known as Run 2, can be found in Table 3.1. Figure 3.3 shows the integrated luminosity over Run 2 as well as the $\langle\mu\rangle$ profile for each of these years.

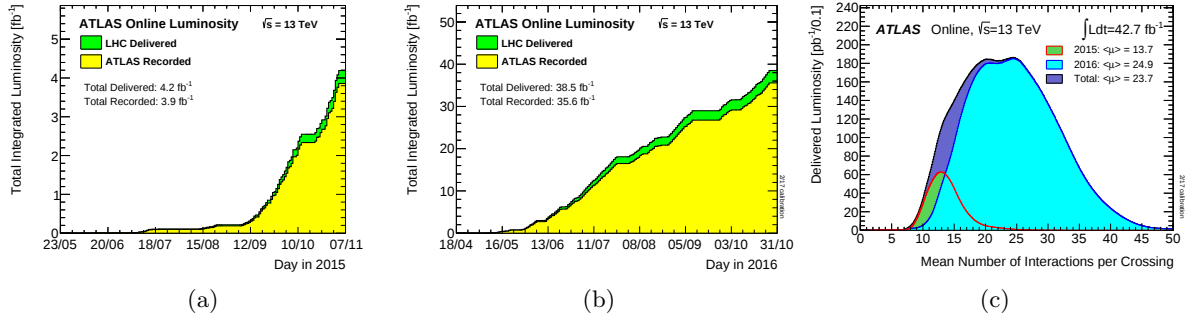


Figure 3.3: The integrated luminosity extracted by the ATLAS experiment during a) 2015 and b) 2016 operations as well as (c) the $\langle\mu\rangle$ profile of these runs [87]. Only 36.1 fb⁻¹ of 2015+2016 data was used for physics studies as detailed in Section 3.2.6.

3.2 The ATLAS Detector

A Toroidal LHC Apparatus (ATLAS) is one of the general purpose detectors of the LHC [79]. It is designed to be able to: measure outgoing particles of all types with good resolution, trigger quickly and efficiently, have full hermetic coverage, and withstand the high radiation levels delivered by the LHC. The detector is cylindrically shaped and symmetric around the beam pipe, approximately 25 m in diameter and 44 m long. The main barrel region is complemented by end-cap regions perpendicular to the beam-pipe to further instrument the forward regions. The detector weighs over 7000 metric tons.

Figure 3.4 shows the layout of the ATLAS detector and its sub-detector systems. In the barrel region ATLAS is composed of several layers of detectors, which starting from those closest to the beam-pipe are:

- An Inner Tracking Detector (ID) to measure the position and momenta of charged particles.
- A liquid argon Electromagnetic Calorimeter (ECal) to measure the energy of charged particles. Mainly designed for electron and photon measurements.
- A Tile Hadronic Calorimeter (HCal) to measure the energy of charged and neutral hadrons.
- A Muon Spectrometer (MS) system, to measure the position and momenta of outgoing muons.

Each of these sub-detectors is further detailed in Sections 3.2.1 - 3.2.4. In addition to these sensitive detector elements, ATLAS also implements a series of hardware and software systems for the extraction of only the data of interest, known as triggering, and for monitoring the status of the detector systems, known as data quality. These are detailed respectively in Section 3.2.5 and Section 3.2.6.

The tracking systems require magnetic fields to bend the trajectories of charged particles in order to measure their momenta. This is provided by a solenoid (toroid) magnet system in

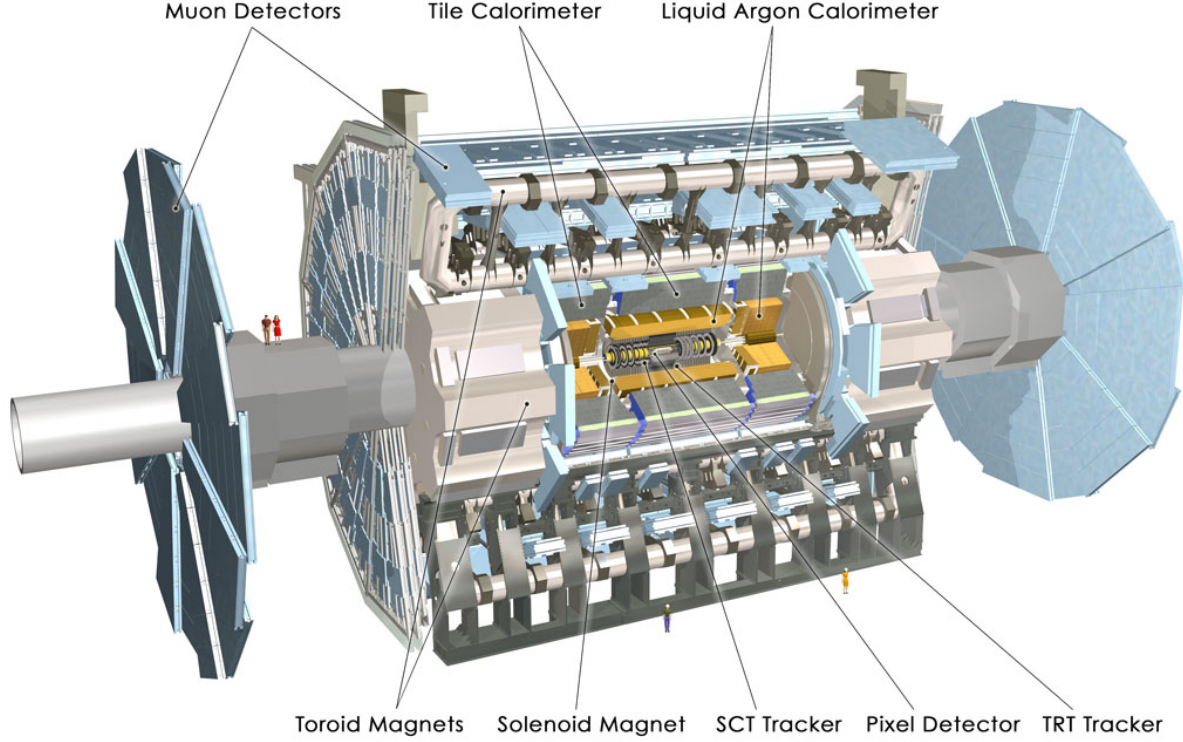


Figure 3.4: Layout of the ATLAS detector [79].

the ID (MS).

The coordinate system used by the ATLAS experiments follows a right-hand system, where the z -axis points along the clockwise rotating beam, the x -axis points towards the center of the ring, and the y -axis points upwards. The interaction point (IP) is defined to be at the origin of this system. The azimuthal angle along the transverse xy -plane is denoted ϕ , and the polar angle above the positive z -axis is denoted θ . The radial distance in the xy -plane is denoted r . The transverse components of observables, such as the transverse momentum p_T , are frequently used and represent the projections along the xy -plane. The pseudo-rapidity is defined as $\eta = -\ln \tan(\theta/2)$, which is related to the rapidity of a particle $y = 1/2 \ln[(E + p_z)/(E - p_z)]$ in the massless limit. The angular separation of two coordinates is defined as $\Delta R = \sqrt{(\Delta\phi)^2 + (\Delta\eta)^2}$.

3.2.1 Inner Tracking Detector

For massive charged particles (those heavier than the electron), the dominant interaction as they pass through matter is through ionization. In this process, the heavy particles ionize atomic electrons, releasing them from the atom and reducing the energy of the charged particle. The mean energy loss due to ionization for a charged particle (charge z) with speed $\beta = \frac{v}{c}$ is

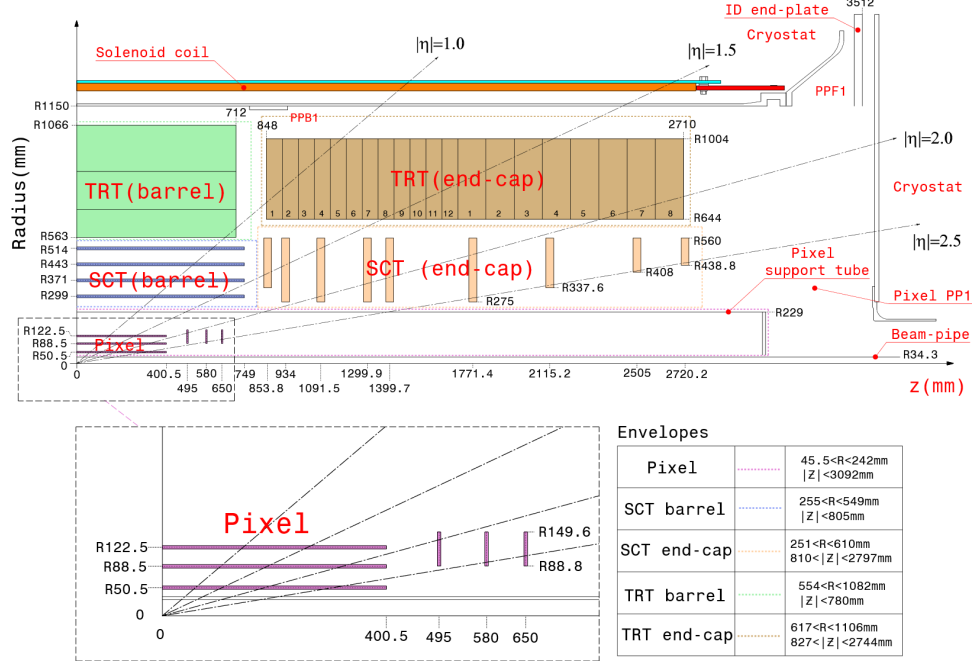


Figure 3.5: Schematic layout of the ATLAS Inner Detector System in the $r - z$ plane [79]. The Insertable B-Layer is not shown.

given by the Bethe-Bloch formula [88].

$$\left\langle \frac{dE}{dx} \right\rangle = -4\pi \frac{nz^2}{m_e c^2 \beta^2} \left(\frac{e^2}{4\pi\epsilon_0} \right)^2 \left(\frac{1}{2} \ln \left(\frac{2m_e \gamma^2 \beta^2 T_{max}}{I^2} \right) - \beta^2 - \frac{\delta}{2} - \frac{C}{Z} \right) \quad (3.7)$$

where n is the average electron density, I is the mean excitation potential, and T_{max} is the max possible energy transfer per collision. The δ and C terms are corrections for density and shell effects respectively. The other constants are the electron charge e , electron mass m_e , and vacuum permittivity ϵ_0 . While the Bethe-Bloch formula provides the average ionization energy, the probability distribution of the energy loss in a finite length absorber is given by the Landau distribution [89]. Qualitatively, the Landau distribution is asymmetric with long tails for high energy losses, which indicate that the most probable energy loss is smaller than the value given by the Bethe-Bloch formula, but high energy deposits are not infrequent.

A tracking system measures the ionizing radiation produced by charged particles as they transverse through the detector. With several layers of sensitive tracking elements the full trajectory of the charged particle can be reconstructed. By applying a magnetic field, the particle momenta can be extracted from the radius of curvature of this path. The sensitive elements are chosen to be sufficiently thick so that the ionization energy can be practically measured, while also being thin enough so that the ionization is not large enough to deter the reconstruction of the charged particle path.

The pp collisions at the ATLAS detector produce thousands of charged particles every bunch-crossing. In order to be able to measure the position and momenta of all of these particles, a

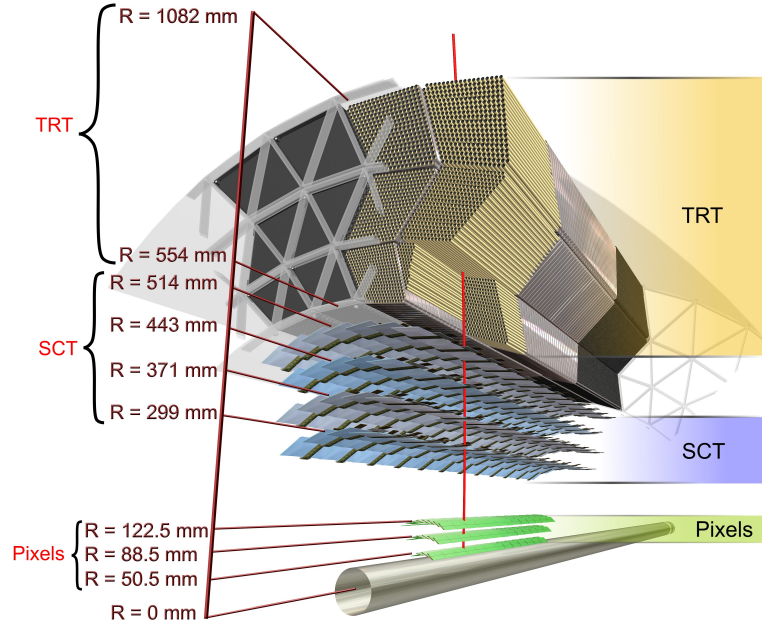


Figure 3.6: Pictorial layout of the ATLAS Inner Detector System highlighting the $r - \phi$ plane in the barrel [79]. The Insertable B-Layer is not shown.

precise and finely instrumented tracking system is needed. In addition, the inner layers must be sufficiently radiation hard as to not degrade too significantly during long-term operation. The Inner Tracking Detector (ID) is composed of a series of Pixel detectors, a Silicon Microstrip Tracker (SCT), and a Transition Radiation Tracker (TRT) composed of straw tubes [90, 91]. The layout of the ID can be seen in Figures 3.5 - 3.6.

The ID extends outwards from just past the beam pipe to a radius of 1.15 m, which is surrounded by a 2 T central solenoid at a radius of 1.23 - 1.28 m of length 5.8 m [92]. The solenoid magnet is composed of a single coil of Al-stabilized NbTi conductor wound 1154 times, in which 7.73 kA of current is passed through. The direction of the magnetic field provided by the central solenoid is along the z direction.

This design allows for excellent track and vertex reconstruction within $|\eta| < 2.5$ for charged particles with $p_T > 500$ MeV. The design momentum resolution expected of the ID is

$$\frac{\sigma(p_T)}{p_T} = 0.05\% p_T \oplus 1\%$$

where $\oplus$ indicates the sum in quadrature of the terms. The individual subsystems are further discussed in the following sections.

Pixel Tracking Detector

The main Pixel system covers an area within $|\eta| < 2.5$, with three layers of silicon pixel detectors centered concentrically in the barrel region at radii of 50.5, 88.5, and 122.5 mm, each 801 mm long in z [93]. In the end-cap region the pixels are placed in disks perpendicular to the beam-pipe at distances of $|z| = 495, 580, \text{ and } 650$ mm. Each pixel has a size of $50 \times 400 \mu\text{m}^2$ in the $r - \phi$ and $250 \mu\text{m}$ thick, which provides measurement accuracy of $10 \mu\text{m}$ in $r - \phi$ and $115 \mu\text{m}$ in $z(r)$ for the barrel (end-cap) region.

During the intermediate period between Run 1 and Run 2 an additional pixel layer, called the Insertable B-Layer (IBL), was installed [94, 95]. The IBL was designed to have improved radiation hardness with respect to the original Pixel system, which was already deteriorating during the conditions of Run 1 which were exceeded in Run 2. The addition of an extra pixel layer closer to the beam-pipe also improves physics performance with better identification of b -hadrons (see Section 4.3.2). The IBL is designed similarly to the Pixel barrel system and placed approximately at $r = 33.25$ mm, and spanning $|z| < 330$ mm. The IBL pixels are smaller than the rest of the Pixel system, with a surface area of $50 \times 250 \mu\text{m}^2$.

The Pixel system geometry guarantees that a charged track should pass through at least four layers of tracking elements within its coverage.

Silicon Microstrip Tracker (SCT)

The next layer of the ID system is a series of silicon microstrip detectors, each of dimension $63.6 \times 64.0 \text{ mm}^2$ with a strip pitch of $80 \mu\text{m}$ [96, 97]. An individual SCT layer only provides measurements in the plane of the strip readouts. To provide unambiguous tracking, a single module of the SCT is comprised of two layers of strips at a stereo-graphic angle of 40 mrad . The SCT system is instrumented to provide the same $|\eta| < 2.5$ coverage as the Pixel system. In the barrel there are 4 layers of SCT modules at radii of 299, 371, 443, and 514 mm, each 1.49 m long. In the end-cap region there are 9 SCT layers at various radii and spatial extent. The end-cap modules are oriented along disks similar to the Pixel system, with one strip layer oriented radially.

As a charged particle passes through the SCT the position is measured 4 times as it passes through eight strip layers. The SCT design provides a measurement accuracy of $17 \mu\text{m}$ in $r - \phi$ and $580 \mu\text{m}$ in $z(r)$ for the barrel (end-cap) region.

Transition Radiation Tracker (TRT)

The last layer of the ID is a collection of 4 mm diameter straw drift-tubes filled with a gas mixture of 70% Xe, 27% CO_2 and 3% O_2 [98, 99]. In the barrel region the straws are 1.44 m long and oriented parallel to the beam-pipe. These straws are divided in two at $\eta = 0$. The TRT barrel extends a radial range of 0.56 - 1.06 m, leading to coverage in $|\eta| < 1$. The modules in this region are segmented into 32 quadrilateral prism layers in ϕ and 3 in r , with 329, 520,

and 793 straws per prism in increasing r layer. The end-caps are composed of straws of length 370 mm, orientated radially into wheels. Each wheel is composed of eight planes of straws, 768 radial straws per plane. The spacing between the planes differs over the wheels for an approximately uniform number of hits in $1 < |\eta| < 2$.

Each barrel (end-cap) straw tube only provides information in the $R - \phi$ ($z - \phi$) plane, with a drift-time measurement allowing accuracy of $170 \mu\text{m}$ per straw. The TRT design allows for typically 36 measurements of the track position for tracks with $p_T > 0.5 \text{ GeV}$ and are within $|\eta| < 2$. Even though the TRT is coarser than the Pixel and SCT system, the large radial extend and high measurement multiplicity of the TRT provides roughly an equal contribution to the track momentum resolution as the other ID layers.

Highly relativistic particles ($\gamma > 1000$) which pass through the polymer fibers in the spaces between the straws can produce photons known as transition radiation. Low energy transition radiation can be absorbed by the Xe-based gas mixture which leads to higher signal values than those of the usual ionization signal. On a straw-by-straw basis, separate low and high-threshold readouts are used. This information is used for electron identification as described in Section 4.3.3. After 2012 operations several Xe leaks were found and some straws were filled with a less expensive Ar mixture instead.

3.2.2 Calorimeters

The interactions of electrons and photons with matter are more complicated than the ionization formula of Equation 3.7. In addition to ionization¹, electrons can also undergo bremsstrahlung with heavy nuclei. The energy loss per length of the material due to bremsstrahlung for a charged particle with mass m and charge z follows

$$\frac{dE}{dx} = 4\alpha N_A \frac{z^2 Z^2}{A} \left(\frac{1}{4\pi\epsilon_0} \frac{e^2}{mc^2} \right)^2 E \ln \left(183 Z^{-\frac{1}{3}} \right) \quad (3.8)$$

where Z and A are the atomic and mass number of the material, and N_A is Avogadro's number. Equation 3.8 scales with E in comparison to $\ln E$ in Equation 3.7 but is suppressed by m^{-2} . This implies that for very high energies, bremsstrahlung is the dominant effect for electrons, which produces a high energy photon. The typical interaction of high energy photons ($>10 \text{ MeV}$ in lead [88]) with matter is for the photon to decay into electron-positron pairs². These two processes continuously chain together, producing showers of lower energy particles until the critical energy E_{crit} is reached, which defines the region where ionization and bremsstrahlung effects are comparable. The longitudinal length scale of these electromagnetic showers is typically given in terms of radiation lengths X_0 , which give the distance for an electron to have only $1/e$ of its initial energy, or equivalently $\frac{7}{9}$ of a photon's mean free distance. The radiation length and

¹A modified version of Equation 3.7 is needed to describe the ionization interaction of electrons and positrons through matter.

²At lower energies, the photoelectric effect and Compton scattering are dominant.

critical energy is material dependent. For copper, $X_0 = 1.43$ cm and $E_{crit} \approx 20$ MeV [88]. The longitudinal length of a shower initiated by an electron or photon of energy E_0 is approximated by $X = X_0 \ln_2(E_0/E_{crit})$. The transverse shower width is given by the Moliere Radius R_M which is the cone width that contained 90% of the shower's energy and can be approximated by $R_M = X_0 \frac{E_s}{E_{crit}}$, where $E_s \approx 21$ MeV.

A similar process, but through different mechanisms also occurs for hadronic particles interacting with dense nuclei. The main process is the production of lower mass hadrons and possible nuclear resonances during the collision of high energy hadrons with a nucleon. These secondary particles can undergo similar interactions until they are less energetic than the pion production threshold, or experience radioactive decays, producing further tertiary particles. These hadrons can also produce ionizing radiation or in the case of π^0 or η mesons, decay to pairs of photons initiating electromagnetic showers. A significant fraction of the energy is also unobserved due to outgoing neutrinos, nuclear fragments such as slow neutrons, or meta-stable particles such as K_L^0 mesons and muons. The longitudinal hadronic shower length in a material is characterized by mean free distance, known as the interaction length, λ_I , which scales approximately as $A^{\frac{1}{3}}$. The number of particles in hadronic showers is lower than in electromagnetic showers with significant variations in the energy sharing between particles.

A calorimeter is a detector designed to both provide material to induce such showers and have the ability to measure the energy deposited during the shower development. In order to accurately measure both the energy of electromagnetic and hadronic particles in a wide geometric range, a series of sampling calorimeters each using different technologies is used [100]. In the barrel region a set of lead-LAr (liquid argon) and steel-tile scintillating calorimeters are used to measure the electromagnetic and hadronic components of the showers. The corresponding end-cap regions and an additional forward calorimeter all use LAr technology, with varying absorber material. Each of these calorimeters is “non-compensating”, which indicates that they do not have equal energy response to hadronic and electromagnetic showers. This effect is accounted for during the calibration of hadronically decaying objects detailed in Section 4.3.2.

These calorimeters are designed such that in total they provide coverage up to $|\eta| < 4.9$, with fine granularity and good energy resolution. In addition, the calorimeter system needs to sufficiently contain the showers so that they do not punch-through into the MS system. The electromagnetic calorimeters provide greater than 22 (24) X_0 radiation lengths of material in the barrel (end-cap) regions. Hadronic showers are similarly contained by approximately 9.7 (10.0) λ_I interaction lengths of material in the barrel (end-caps). This level of shower containment provide reliable estimates of the jet energy and missing transverse energy (described in Section 4.3.2 and Section 4.3.5 respectively) over a wide energy range. Each of the individual ATLAS calorimeters is discussed in further detail in the subsequent sections.

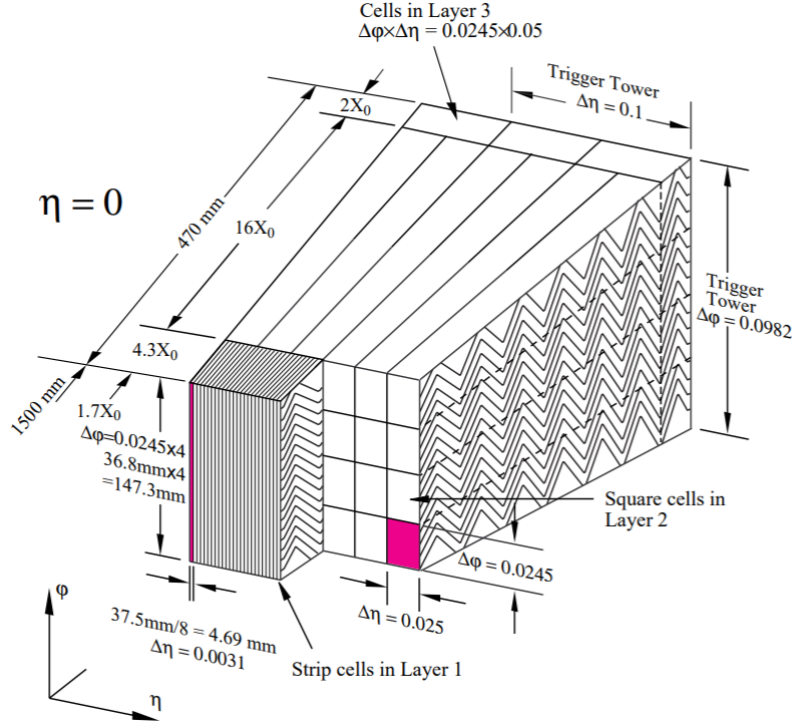


Figure 3.7: Diagram of barrel ECal design and readout granularity [79].

LAr Electromagnetic Calorimeter

The barrel electromagnetic calorimeter (ECal) and electromagnetic end-cap (EMEC) use accordion shaped absorbers/electrodes using lead as the active absorber material and liquid-argon (LAr) as the sampling material [101]. The accordion structure, as seen in Figure 3.7 extends along $r(z)$ alternating in ϕ for the ECal (EMEC), allowing for almost uniform coverage in ϕ without cracks. The lead thickness varies as a function of η to optimize resolution performance, ranging from 1.1 - 2.2 mm. In the barrel, the total gap width is 4.2 mm, but varies in the end-cap as the accordion waves vary with radius. At 2000 V, the signal pulse length is 450 ns.

The inner and outer radii of the barrel calorimeter are 1.4 and 2 m respectively and both half barrels cover a $|z| < 3.2$ m. The end-caps are composed of two wheels, 0.63 m thick, extending from $r = 0.33$ to 2.10 m. This provides a coverage of $|\eta| < 1.475$ for the ECal and $1.375 < |\eta| < 3.2$ for the EMEC. The barrel ECal is split into two at $\eta = 0$, with a total of 1024 accordion absorbers broken up into 16 modules spanning $\Delta\phi = 22.5^\circ$. Each module has three longitudinal readout segments, 4.3, 16, and 2 X_0 long respectively. The $\Delta\eta \times \Delta\phi$ readout granularity increases per longitudinal layer and varies over η . The granularity of the first (third) barrel layer is $0.025/8 \times 0.1$ (0.05×0.25) in $\Delta\eta \times \Delta\phi$. Figure 3.7 displays the readout granularity of the ECal modules. The EMEC is separated into inner, $1.375 < |\eta| < 2.5$, and outer, $2.5 < |\eta| < 3.2$ wheel, each containing 8 azimuthal modules. In the region $1.5 < |\eta| < 2.5$,

the modules are divided into 3 longitudinal layers, similar to the barrel. The outermost regions of the EMEC $|\eta| < 1.5$ and $2.5 < |\eta| < 3.2$ are divided into only two longitudinal sampling layers. The $\Delta\eta \times \Delta\phi$ granularity varies per module and layer with the finest and coarsest matching those of the barrel region.

In order to improve the energy measurement in $|\eta| < 1.8$, where there is significant dead material before the calorimeter, an additional pre-sampler calorimeter layer is placed. The pre-sampler layers are single 11 mm layers of LAr in the barrel and two 2 mm layers of in the end-cap. They are divided to have granularity of 0.025×0.1 in $\Delta\eta \times \Delta\phi$.

The design specification of the energy resolution of the ECal

$$\frac{\Delta(E)}{E} = \frac{10\%}{\sqrt{E}} \oplus 0.7\% \quad (3.9)$$

Tile Hadronic Calorimeter

In the barrel region, a steel-tile scintillator sampling hadronic calorimeter (HCal) is used for the containment and measurement of hadronic components of showers [102]. The HCal surrounds the ECal from $r = 2.28$ to 4.25 m. It is broken up into a central barrel 5.64 m long in z and two extended barrels 2.91 m long, which together cover $|\eta| < 1.6$. Each module, which divides the HCal into 64 azimuthal section follows a periodic structure of tiles of steel absorber 14 mm thick and sampling scintillator 3 mm thick. The wedge-shaped tiles are oriented such that the largest face of the tiles span $\phi - r$. Figure 3.8 shows the orientation of tiles within a HCal module

The light from the scintillating tiles is readout by wavelength-shifting fibers, which are grouped and fed into photo-multiplier tubes. The fiber grouping defines the three-dimensional readout granularity. Radially it is segmented into three layers of approximately 1.5, 4.1, and $1.8 \lambda_I$ in the barrel, and 1.5, 2.6, and $3.3 \lambda_I$ for the extended barrel. The ECal provides approximately $2 \lambda_I$ of material as well. The grouping in η provides the first two layers with a granularity of 0.1×0.1 in $\Delta\eta \times \Delta\phi$ and 0.2×0.1 in the final layer. The physical readout grouping for each HCal module can be seen in Figure 3.9.

The energy resolution of the HCal during design was

$$\frac{\Delta(E)}{E} = \frac{50\%}{\sqrt{E}} \oplus 3\% \quad (3.10)$$

Hadronic End-cap Calorimeter

In the region of $1.5 < |\eta| < 3.2$ a hadronic end-cap calorimeter (HEC) using LAr technology similar to the ECal is used [101]. In comparison to the ECal, the HEC uses copper as the absorber material and a flat-plate design as opposed to the accordion structure. The layout of the HEC, shown in Figure 3.10, is in the form of a front and back wheel, 1.82 m in depth. The first wheel is composed of 24 copper plates 25 mm thick, while the second wheel is composed

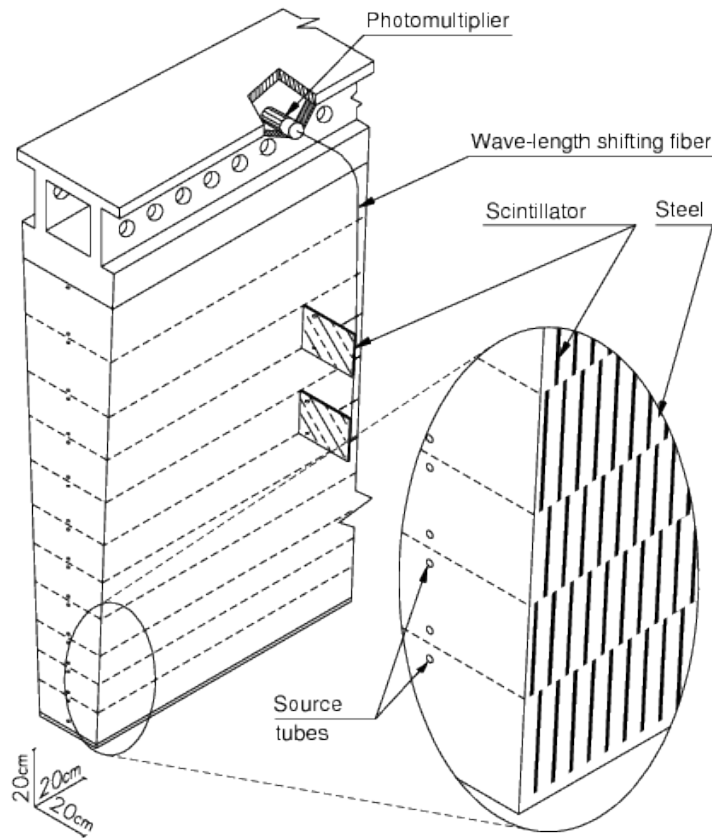


Figure 3.8: Diagram of one HCal module and the orientation of tiles in the module [79].

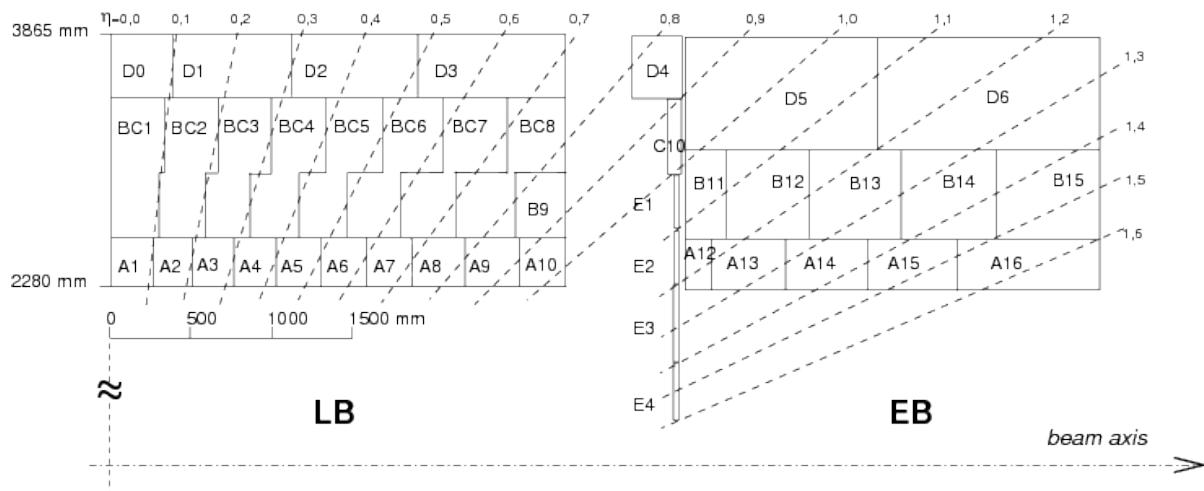


Figure 3.9: Segmentation of the HCal readouts in η and radius for the barrel (left) and the extended barrel (right) [79].

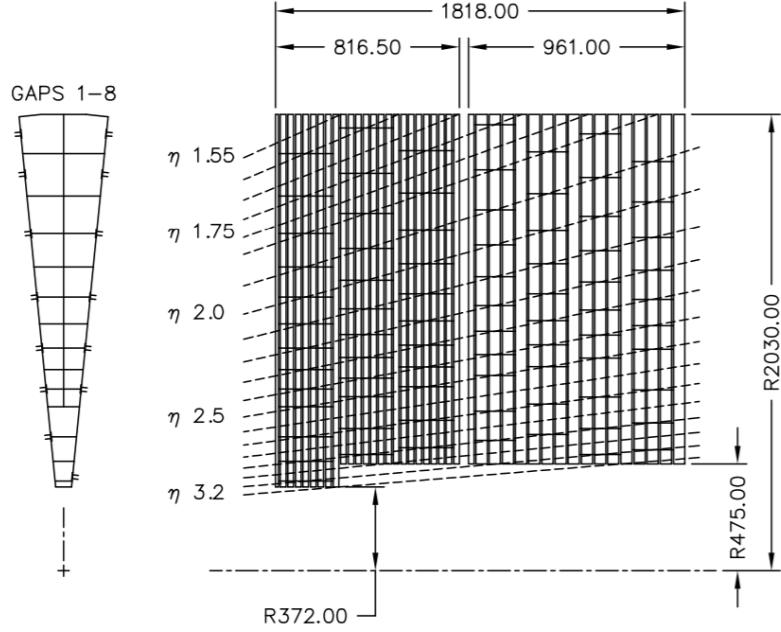


Figure 3.10: Views of the HEC module division in the $r - \phi$ plane (left) and $r - z$ plane (right) [79].

of 16 plates 50 mm thick. The gaps between each plate are 8.5 mm wide and filled with LAr as a sampling material. The gaps are instrumented with three planar electrodes, instrumented such that the readout can be treated as two double gaps, allowing for lower applied voltages. For an applied voltage of 1800 V this leads to an expected signal pulse 430 ns wide.

The wheels are divided into 32 azimuthal sections, and twice along z for each of the two wheels. This leads to 4 sampling layers laterally, but with varying granularity in η . Within $1.5 < |\eta| < 2.5$ the granularity in $\Delta\eta \times \Delta\phi \sim 0.1 \times 0.1$, while in $2.5 < |\eta| < 3.2$ it is approximately 0.2×0.2 . The energy resolution of the HEC is similar to the barrel HCal.

Forward Calorimeter

Calorimetry in the far forward region ($3.1 < |\eta| < 4.9$) requires dedicated detector technology to withstand the high amounts of radiation from the high particle flux in this rapidity region. In addition, the physical volume in this region is quite small and a dense absorber is needed to contain the showers and protect the other ATLAS detectors. This is done with a series of LAr based detectors known as the Forward Calorimeter (FCal) [101]. The FCal is composed of thin < 0.5 mm electrode gap rods segmented into three layers along z . The rods run along the direction of the beam-pipe (z) and are 0.45 m long per layer. Figure 3.11b shows the position of the FCal layers in the context of the ATLAS beam-pipe and end-cap calorimeters.

The first FCal layer acts as an electromagnetic calorimeter, with copper acting as the absorber material. The modules here are composed of copper rods with 0.27 mm LAr gaps

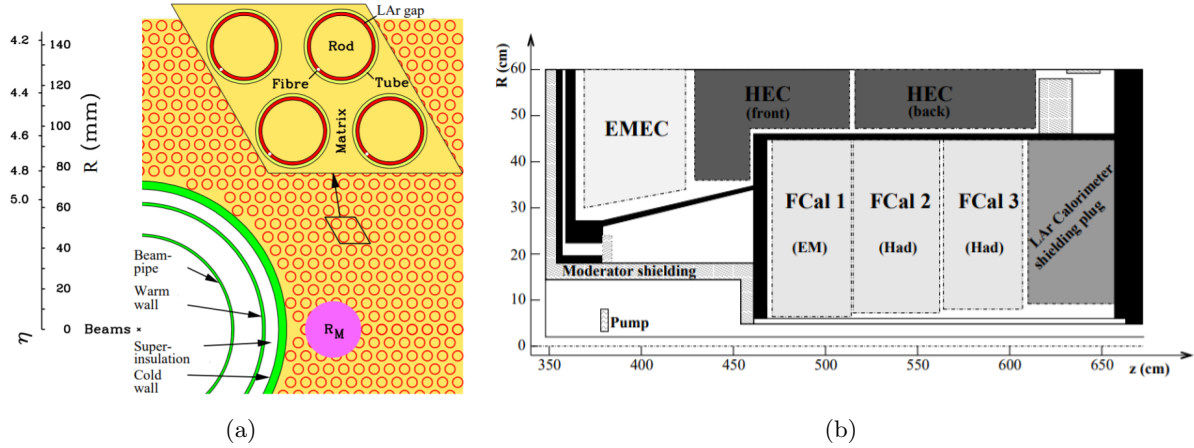


Figure 3.11: a) Diagram of the cross-section of the first FCal layer. The size of the expected Moliere radius is also shown. b) The layout of the ATLAS detector calorimeters around the FCal region [79].

between an outer copper tube with radius 2.5 mm. This leads to a 60 ns drift time. These tubes are placed into holes drilled into a copper bulk in a hexagonal matrix pattern, with 7.5 cm between the centers of the rods. The tube and matrix structure of the first FCal layer can be seen in Figure 3.11a. The next two FCal layer acts as hadronic calorimeters with tungsten as its main absorber material. The structure is similar to the first FCal layer but the inner rods and bulk material are now compose of tungsten as opposed to copper. The LAr gap is also increased to 0.37 (0.51) mm in second (third) layer.

In total, the FCal provides approximately $10 \lambda_I$ and $208 X_0$ of material for shower containment. The design energy resolution of the FCal is

$$\frac{\Delta(E)}{E} = \frac{100\%}{\sqrt{E}} \oplus 10\% \quad (3.11)$$

3.2.3 Muon Spectrometer

The final layer of the ATLAS detector is a tracking detector designed to measure high energy muons, called the Muon Spectrometer (MS) [103]. Due to the small amount of ionizing radiation and bremsstrahlung expected from muons, they typically are not expected to be fully contained by the calorimeter systems. For example, it is expected that a 100 GeV muon will lose approximately 3.8 GeV over $100 X_0$ [88]. To measure such particles, the MS is a tracking system composed of precision tracking layers and dedicated trigger systems.

In a similar approach to the ID, measurement of the particle momenta can be inferred from particle trajectory in a magnetic field. The magnetic field is provided by a series of superconducting air-core toroid magnets [104, 105]. The barrel toroid magnet covers the region of $|\eta| < 1.6$, while two end-cap toroids cover $1.4 < |\eta| < 2.7$. Each toroid system is composed of eight flat race-track shaped coils spanning $z - r$, and spaced evenly in ϕ . The magnetic field

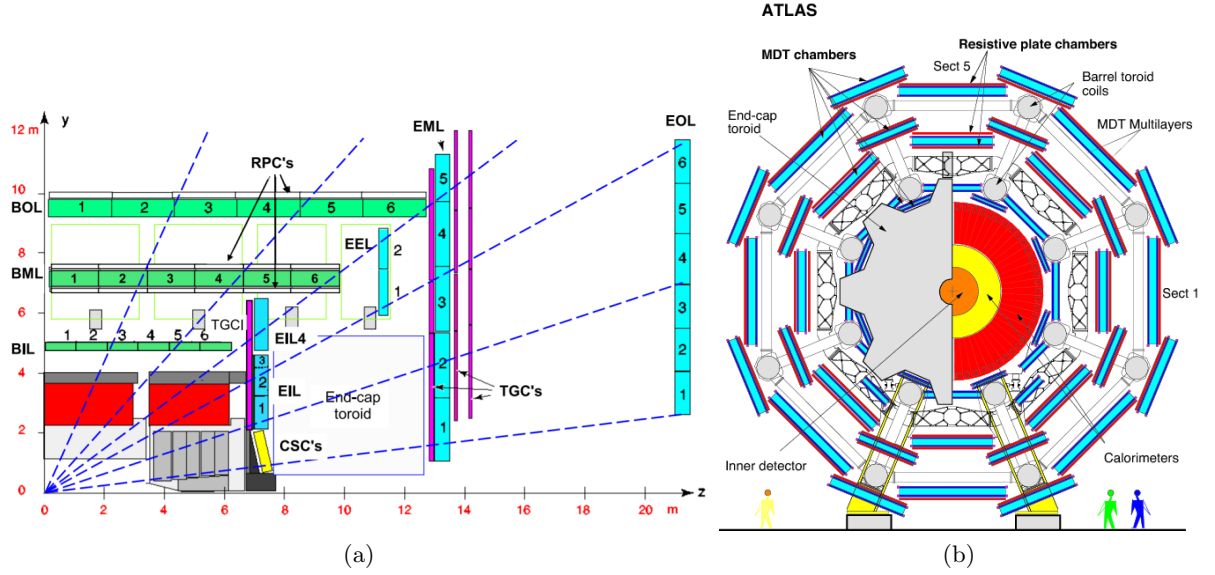


Figure 3.12: Schematic layout of the Muon Spectrometer system in a) the $r - z$ plane [79] and b) the $r - \phi$ plane [106].

can vary over the large spatial extend of the toroid magnets, providing an integrated magnetic field of 1.5 - 5.5 (1.0 - 7.5) Tm in the barrel (end-cap) regions.

The MS tracking system is comprised of a Monitored Drift Tubes (MDT) which covers most of the η range and Cathode Strip Chambers (SCT) in the end-cap. The trigger systems, which also provide measurements along the orthogonal direction with respect to the MDT/SCT readout, are similarly divided into Resistive Plate Chambers (RPC) and Thin Gap Chambers (TGC), in the barrel and end-cap. The eightfold symmetry of the toroid structures are mimicked also in the design of the tracking and triggering systems. Figure 3.12 shows a schematic layout of the MS systems. Each system is detailed in subsequent sub-sections.

The design momentum resolution of the MS is [103]:

$$\frac{\Delta(p_T)}{p_T} = 0.01 p_T \quad \text{for } p_T > 300 \text{ GeV} \quad (3.12)$$

This provides the common benchmark value of a 10% uncertainty on 1 TeV muons. This corresponds to a precision of 30 μm on the relative position of the tracking elements. To maintain this precision over the large detector volume of the MS, an optical high-precision alignment system is used to measure the positions of all the tracking elements. The muon resolution, especially at low momenta, can be improved by combining with ID tracks as detailed in Section 4.3.4.

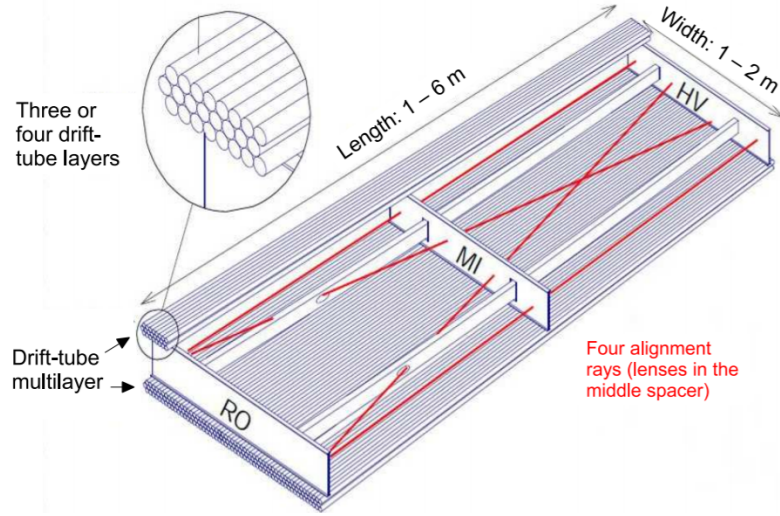


Figure 3.13: Schematic layout of one of the MDT chambers [79]. The two layers of drift tubes can be seen along with the optical positioning system.

Monitored Drift Tubes

The main tracking system of the MS is a series of Monitored Drift Tubes (MDT) [103]. The MDT is composed of a series of 30 mm diameter drift tubes filled with a 97:3 ratio of Ar:CO₂ and a tungsten-rhenium cathode wire at approximately 3000 V. In the barrel, the MDT system contains three planar chambers, approximately at $r = 5, 7.5, 10$ m in a eightfold concentric geometry, mounted either between the barrel toroid gaps, or on the barrel toroid support as seen in Figure 3.12. In the end-cap, they are layered perpendicular to the beam-pipe at $|z| = 7.4, 10.8, 14,$ and 21.5 m with spatial extent such that straight-tracks must pass through at least 3 layers in $|\eta| < 2$. In the range $2 < |\eta| < 2.7$, the MDT only covers the outer two layers while the CSC covers the inner layer.

Each of the MDT chambers is composed of two groups of tube layers, with the tubes oriented along the ϕ direction. In the innermost chambers there are four tubes per layer, while in the outer there are three. The geometry of each chamber varies per module to provide the maximum solid angle coverage. The gaps between the chamber layers are instrumented with an optical ray system which can measure deformations to within $30 \mu\text{m}$. Since the tube layout is continuously measured and corrected for, the dominant position uncertainty is the intrinsic tube resolution of $80 \mu\text{m}$, or $35 \mu\text{m}$ per chamber. Figure 3.13

During the shutdown preceding Run 2, some of the missing MDT chambers in the transition region $1.0 < |\eta| < 1.4$ were added completing the above design. These new MDT chamber have a smaller tube radius then the rest of the MDT system.

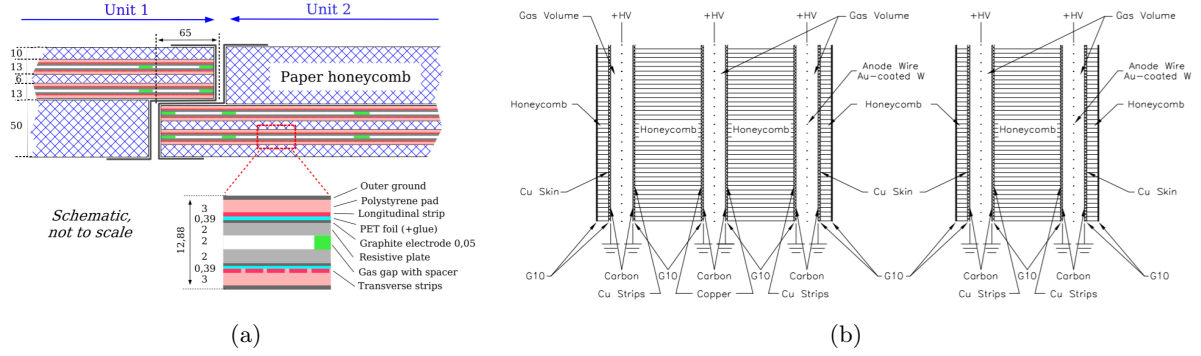


Figure 3.14: Structure of a) one RPC doublet b) the TGC doublet and triplet chambers [79].

Cathode-Strip Chamber

As mentioned in the previous section, the inner-most layer of the MS tracking system in $2 < |\eta| < 2.7$ is made of Cathode-Strip Chambers (CSC) [103]. The different detector technology is required in this region to withstand the higher particle flux.

The layout of the CSC is similar to the MDT in the end-cap, with two wheels segmented eightfold in ϕ . The CSCs are multi-wire proportional chamber with the wires oriented radially down the wheels and cathode walls segmented perpendicular or along the wire direction for readout in both directions. Each chamber is filled with a 80:20 Ar:CO₂ mixture with a gas gain of 6×10^4 for an operating voltage of 1.9 kV. The anode-cathode distance and anode pitch are both equal at 2.5 mm. Each CSC layer provides a space point, and each chamber is made of four CSC layers.

The readout, composed of three individual cathode strips, for the two wheels have a pitch along the bending plane of 5.3 and 5.6 mm, and 13 and 21 mm in the non-bending plane. Signal pulses along the three to five most-neighboring cathode strips are interpolated to provide further precision. This leads to 60 μm resolution in the bending plane and 5 mm in the other direction.

Resistive-Plate Chamber

To complement the high-precision tracking, fast trigger chambers are used. In the $|\eta| < 1.05$ region these are provided by Resistive Plate Chambers (RPC) [103]. The RPC is constructed by two parallel resistive plastic laminate plates, separated by 2 mm. The gaps between the plates are filled with a C₂H₂F₄/Iso – C₄H₁₀/SF gas mixture, with an applied electric field of 4.9 kV/mm. Incident particles can ionize the gas and cause avalanches in the gas mixture, which are read out via capacitive coupling of metallic strips on the outside of the plates. The readouts have a pitch 23 - 35 mm in η and ϕ , varying per chamber with signal widths of 5 ns.

Each RPC trigger chamber is made of two doublets of RPCs, providing two sets of η and ϕ measurements per chamber as seen in Figure 3.14a. Two RPC chambers are placed on both sides of the middle MDT layer and a final is placed on the outside of the third MDT layer. This

structure allows triggering on both high and low p_T muons due to the different bending arms.

Thin-Gap Chambers

In the region $1.05 < |\eta| < 2.4$ a finer granularity trigger system is needed due to higher radiation, lower integrated magnetic fields, and the requirement for uniform triggering in p_T . Thin-Gap Chambers (TGC) are used in this region to provide the necessary finer readout [103]. The TGC are multi-wire proportional chambers similar to the CSC. The gas is 55:45 $\text{CO}_2/\text{n-pentane}$ mixture with a gas gain of 3×10^5 at an operating voltage of 2.9 kV. The anode-cathode distance is 1.4 mm while the anode pitch is 1.8 mm. This design allows a more uniform performance of pulse-height with respect to incident angle and provides 15 ns pulse widths.

The TGC are constructed in modules of doublet or triplet chambers as seen in Figure 3.14b. Each TGC chamber has two graphite cathode plates, on which copper layers are attached. Two of these copper layers per doublet or triplet are segmented for readout in the azimuthal coordinate while the other serve as grounding. The measurement of the radial bending coordinate is provided by anode wires. To provide the required granularity in all η regions, between 4 and 31 anode wires are grouped and readout together. The inner wheel of the MDT end-cap has one TGC doublet, while the middle wheel has one triplet and two doublets. The match of MDT-TGC hits for the last MDT layer is extrapolated from the last TGC layer, which is past the end-cap toroid field.

3.2.4 Forward Detectors

The ATLAS experiment is also instrumented with a set of far-forward calorimeters, which measure the particle flux very close to the beam directions at large $|\eta|$. These include the LUCID2 detector (simplified here to LUCID) [107] and the Beam Conditions Monitor (BCM) [108]. Both of these include luminosity determination, as detailed in Section 3.1.1, as one of their physics goals. Both detectors are described in further detail in subsequent sections

LUCID

The LUCID detector is a Cherenkov radiation detector located $z = 16.97$ m from the main ATLAS interaction point [107]. The version of LUCID used in Run 2 is an upgraded version of the original detector, which was significantly radiation damaged during Run 1. This version of LUCID is composed of 16 photomultipliers (PMTs) of 5 mm radius, grouped into fours. Originally in each bundle was one PMT with reduced acceptance of 3.5 mm, one calibrated with LED signals, one calibrated with ^{207}Bi sources, and one spare identical to the LED calibrated PMT. The calibration sources are placed directly in front of the PMTs and allows for calibration of the PMT voltage to provide constant gain to compensate for deterioration during operation. The Bismuth calibrated device showed the best performance and the spare was switched with Bismuth calibrated devices in 2015, and the other two remaining devices in 2016.

The main PMTs and quartz fibers are on aluminum supports just outside the beam-pipe at radius 125.5 mm.

As charged particles pass through the quartz window of the PMT, they produce Cherenkov radiation which is picked up by the PMT itself. The threshold for a LUCID signal in a PMT is 15 photo-electrons, with the mode being 30. LUCID can also measure the integrated pulse and provide luminosity values for each bunch-crossing.

During Run 2 operation several PMTs stopped working. By the end of 2018 only a single PMT was used for measurement of luminosity.

Beam Conditions Monitor

The Beam Conditions Monitor (BCM) is a dual purpose detector designed to measure both the bunch-by-bunch luminosity as well as trigger beam aborts in the situation where the beam is mis-aligned [108]. In the situation where the proton beams becomes greatly mis-aligned and hits one of the collimators, large amounts of radiation can be released, and while the ATLAS detectors are designed to withstand high radiation doses, many would be greatly damaged by the high particle fluxes produced in such events. To identify such possible occurrences, the BCM is composed of four $8 \times 8 \text{ mm}^2$ diamond pixel pads, $500 \text{ }\mu\text{m}$ thick, in a cross-pattern around the beam-pipe, $|z| = 1.84 \text{ m}$ away on each side of the interaction point. Each module is positioned just beyond the beam-pipe at $r = 55 \text{ mm}$, and hence at $\eta \approx 4.2$. Chemical vapour deposited diamond is used as it has a large electron band-gap of 5.5 eV (c.f. 1.1 eV in Si) which makes it intrinsically radiation hard [109].

Ionizing particles will leave on average 36 electron-hole pairs per mm as they traverse the BCM modules. A bias voltage of $2 \text{ V}/\mu\text{m}$ is applied across the detectors, which cause the electron-hole pairs to drift apart before being trapped after some distance. In the chemical vapour deposited diamond used by the BCM, this distance is approximately $300 \text{ }\mu\text{m}$ for fields above $0.5 \text{ V}/\mu\text{m}$.

The BCM signal has a 4.5 ns pulse width, allowing bunch-by-bunch measurements. Due to this fast response time and position of the detectors, if coincident hits are measured on both sides of the detector, the BCM can distinguish whether these originate from two particles originating from the interaction point or one particle traveling parallel to the beam-pipe through the detector. For the former, the time difference between the two coincident hits will be approximately 0. For the later, the time difference will be approximately 12.5 ns, within the response time of the BCM. If the rate of out-of-time coincidences becomes too large, the BCM signals a beam abort to protect the other ATLAS detectors.

3.2.5 Trigger and Data Acquisition System

The trigger is an online (during the LHC operations) system which decides whether to record an event or not. Due to hardware constraints on electronic readout speeds and storage sizes, it is impossible to store every LHC bunch-crossing which can occur at a rate of 40 MHz.

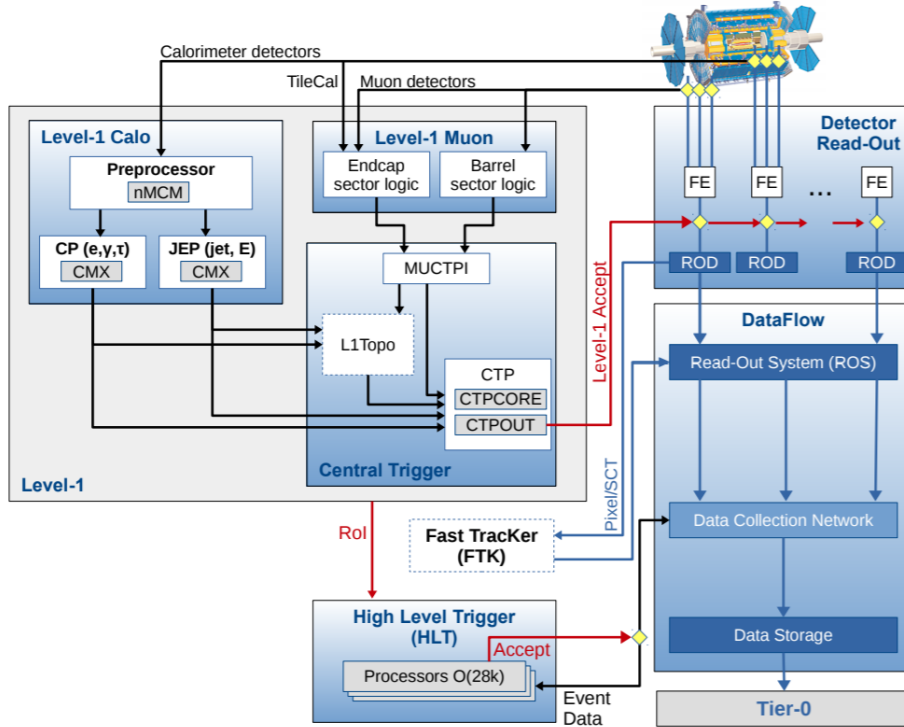


Figure 3.15: Overview of the ATLAS Trigger and Data Acquisition Systems. The Fast Tracker (FTK) system was not used in Run 2 operations.

The majority of these events are “uninteresting” low-energy collisions, so the trigger system is designed to quickly search for signatures of high-energy collisions to be stored for offline (after data taking) analysis.

The trigger system is composed of a Level 1 (L1) [110] hardware based system which reduce the readout rate to 100 kHz, followed by a software based High Level Trigger (HLT) [111] which further reduces it to an 1 kHz rate. The overall latency of the L1 trigger is under $2.5 \mu\text{s}$ and 400 ms for the HLT. A diagram of the data-flow in the Trigger and Data Acquisition System (TDAQ) is shown in Figure 3.15. Both the L1 and HLT have changed significantly from their original design for operations during Run 2. The HLT was originally two separate software systems which have been merged and optimized, while the L1 system has seen extensively upgraded to cope with higher trigger rates [112].

Level 1 Trigger

The Level 1 (L1) trigger makes a decision from a coarse readout of the calorimeter, RPC and TGC detectors. The L1 system defines Regions-of-Interest (RoI) associated to high- p_T objects passing pre-defined selection, which are fed into the HLT to seed further algorithms.

The L1 calorimeter system identifies possible electron, photons, and hadronically decaying τ -leptons via sliding windows of 2×2 trigger towers in the ECal readout. One trigger tower is

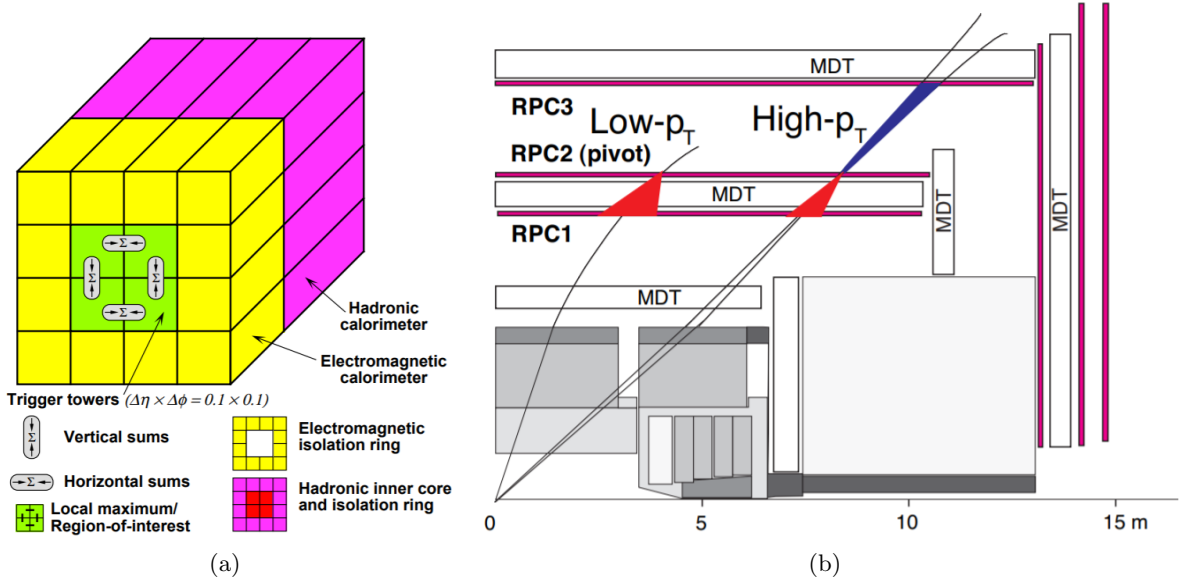


Figure 3.16: a) Diagram of the cross-section of the first FCal layer. The size of the expected Moliere radius is also shown. b) The layout of the ATLAS detector calorimeters around the FCal region. [79].

approximately 0.1×0.1 in $\Delta\eta \times \Delta\phi$, but can increase as a function of η . To pass the trigger one set of neighboring 2×1 or 1×2 trigger towers is required to pass a predefined E_T threshold. To remove overlaps, the window is required to be a local maximum with respect to the surrounding eight windows. Isolation requirements can be set on E_T thresholds in the 16 surrounding trigger towers in the ECal and HCal as well as the 2×2 trigger located in the HCal layer as shown in Figure 3.16a. Jet triggers are identified by 4×4 or 8×8 trigger towers summed between the ECal and HCal, which pass a E_T threshold and the inner 2×2 core is a local maximum.

The L1 muon system operates by attempting to quickly reconstruct tracks through multiple RPC and TGC layers. In the barrel $|\eta| < 1.05$, if the second RPC layer finds a hit, it searches for a corresponding hit in the first layer in a road centered on the straight line from interaction point to the original hit. The road width is based on the p_T threshold used for the trigger. To be considered a valid track, at least three out of the four RPC doublet layers are required to have coincident hits. High p_T triggers are also defined if a road built forward to the outer layers also provides a hit in one of the two RPC doublets. A diagram of the muon road-seeding in the barrel can be seen in Figure 3.16a. The end-cap ($1.05 < |\eta| < 2.5$) muon triggering is done similar for the TGC, with the roads being seeded by the final TGC doublet layer. A 3-out-of-4 coincidence is required between the two outer TGC doublets and 2-out-of-3 in the triplet layer to be considered a valid track. To further suppress background rates in $1.3 < |\eta| < 1.9$ muon triggers require a coincident hit in the innermost TGC layer.

The results and RoI of the L1 muon and calorimeter system are then fed into an L1Topo algorithm which allows further triggering based on geometric quantities (such as relative angles)

or kinematic quantities (such as invariant mass) from several trigger objects as well as a refined E_T^{miss} calculation. All the L1 trigger systems results are then provided to a Central Trigger Processor (CTP) which makes the final L1 decision. A successful pass of the CTP then feeds the RoI information to the HLT as well as indicates for the detector specific readout drivers to read the information from each of the respective electronic front-ends. The data is stored in read-out buffers until the HLT decision is made.

High Level Trigger

After a successful L1 trigger pass, the High Level Trigger (HLT) performs a more precise event reconstruction using precision measurements from the MS, ID, and calorimeters depending on the L1 seed and RoI. The algorithms performed by the HLT are similar to those performed by the offline analysis detailed in Section 4.3.

Jet triggers use the whole calorimeter volume and use an anti- k_t clustering algorithm [113], with radius parameter $R = 0.4$ or $R = 1.0$ on topologically clustered energy deposits [114]. These jets are calibrated to the EM scale and use a similar procedure to the offline analysis (see Section 4.3.2), modulo the *in-situ* measurements, to calibrate to the hadronic scale. Triggers on E_T^{miss} also use calorimeter information but are much more sensitive to energy calibrations. Several E_T^{miss} triggers are defined using cell-based or jet-based calculations with various pile-up subtraction methods.

For muon triggers, the precision measurements from the MDT in the RoI are used to improve the position and p_T measurements of the tracks. For most triggers, there is the additional condition that muon tracks can be back-extrapolated to corresponding ID tracks, or that ID tracks can be forward extrapolated to muon tracks (so called combined muon candidates). ID tracks are found with a HLT specific tracking algorithm which requires at least two pixel hits or six SCT hits in $|\eta| < 2.5$. An isolation requirement is typically also required for most muon triggers which requires the scalar p_T sum of all the tracks within $\Delta R < 0.2$ be less than 12% of the muon candidate p_T .

Electron reconstruction at the HLT is first done with a sliding window algorithm to cluster the finely segmented calorimeter cells by searching for the most energetic cluster in each layer within the RoI. The window size is $\Delta\eta \times \Delta\phi = 0.075 \times 0.175$ in the barrel and 0.125×0.125 in the end-cap, with the cluster center allowed to differ amongst calorimeter layers. A further set of requirements on several shower shape parameters are then applied to remove misidentified hadronic jets (see Ref. [112] for details). To remove photon events, an ID track with $p_T > 1$ and $\Delta\eta < 2$ with respect to the cluster center is also required. Like jets, trigger-level electrons are calibrated separately in the same way as offline algorithms (detailed in Section 4.3.3). Electron candidates are further identified based on likelihood-based discriminant with identical inputs as the offline analysis, but with slightly different distributions as the effects of bremsstrahlung radiation are not included in the HLT.

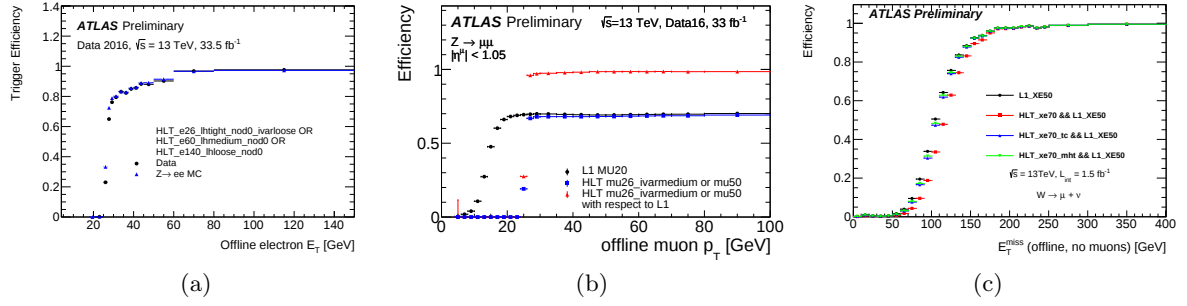


Figure 3.17: The trigger efficiency for the various triggers used in this thesis. a) The efficiency of logical OR of the 2016 electron triggers as measured with tag-and-probe techniques in $Z \rightarrow e\bar{e}$ in data (black) and MC (blue) [115]. b) The efficiency of the logical OR of the 2016 muon triggers as measured in $Z \rightarrow \mu\bar{\mu}$ data (blue) [116]. The L1 seed trigger efficiency (black) and HLT to L1 pass efficiency (red) are also shown. c) The 2015 E_T^{miss} trigger efficiencies as measured using inclusive $W \rightarrow \mu\nu$ events [117]. The black curve shows the L1 seed efficiency and the red shows the efficiency for the trigger used in this analysis

Trigger in Run 2 Operations

The analyses in Chapter 5 and Chapter 7 will make use of the leptonic triggers to select events consistent with a leptonically decaying W/Z boson. In particular, the lowest unprescaled triggers available during that data taking period will be used, which are detailed in this section.

During the 2015 (2016) data taking, the lowest threshold electron triggers require $E_T > 24$ (26) with p_T dependent isolation requirements. In addition, a $E_T > 60$ GeV electron trigger with no isolation requirement and an $E_T > 120$ trigger with looser identification requirements are used. The loosest muon triggers require $p_T > 20$ (26) GeV in 2015 (2016) with isolation requirements applied, or $p_T > 50$ without isolation requirements. These lepton triggers provide approximately 95% and 70% efficiency for events with electrons [115] and muons [116]. For analyses targeting leptonic W decays, such as those in this work, the lower trigger efficiency in muon events can be recovered using E_T^{miss} triggers, since muons are not included in the trigger E_T^{miss} calculation. The lowest threshold E_T^{miss} triggers have thresholds of 70 GeV for the 2015 data and of 90 – 110 GeV for the 2016 data. Tag-and-probe studies found that the E_T^{miss} triggers are fully efficient for events with $p_T(W) > 200$ GeV [117]. The relevant trigger efficiencies as a function of turn-on variable can be found in Figure 3.17.

3.2.6 Data Quality

To guarantee that the data extracted from the ATLAS detector is free of hardware or software related issues, the data is scrutinized before being certified for physics usage in a procedure known as Data Quality (DQ) [118]. This multi-step procedure iteratively stores and updates meta-data on the conditions of the beam, trigger, magnets, and detector elements. During ATLAS operations, a team of “shifters” monitor the data-taking in real-time to identify potential

DQ issues and to resolve them as quickly as possible. A subset of the data is then promptly reconstructed and analyzed by “DQ experts” to identify defects in detector operations and to evaluate the conditions of the data run. With the updated conditions, detector calibrations are adjusted and a processing of the full run data is initiated.

During the DQ procedure possible issues in the detector can be identified, such as noisy SCT strips or LAr calorimeter readouts, which can be masked during the full data reconstruction. More major issues, such as the desynchronization of the pixel readouts or failure of the toroid magnets, are also identified and stored in the conditions meta-data. This later list of issues is parsed to select data runs within detector conditions tolerance, which are designated as “good for physics.” The fraction of recorded data to “good for physics” data is known as the DQ efficiency.

During 2015 - 2016 operations, 42.7 fb^{-1} of pp data was delivered by ATLAS as can be seen in Figure 3.3. Due to various issues which can arise during operations, only 39.5 fb^{-1} of this data was recorded. During the 2015 (2016) ATLAS operations, the DQ efficiency was 88.8% (93.1%) [118]. Thus, the overall 2015 - 2016 ATLAS dataset includes 36.1 fb^{-1} of pp collision data which is “good for physics” and will be analyzed in subsequent chapters.

Summary

The ATLAS detector is designed to hermetically measure the particles originating from the 13 TeV pp collisions from the LHC. Through several layers of tracking detectors it can measure charged particle tracks with p_T above 0.5 GeV and with finely instrumented calorimeters can measure the energy of hadronically and electromagnetically showering objects from 10 GeV up to several TeV. In addition, a dedicated muon system can reconstruct muon candidates with momentum 3 - 3000 GeV. The ATLAS detector also includes a fast triggering system which allows for low-threshold triggering on electron and muon candidates. Dedicated algorithms for reconstructing these particles from the digitized ATLAS signal for physics usage are detailed in Chapter 4.

During the ATLAS Run-2 operations, approximately 36 fb^{-1} of 13 TeV pp collision data was recorded. This data-set will be analyzed in Chapters 5 - 7 to search for new physics signatures in final states with W/Z pairs.

Chapter 4

Detector Simulation and Event Reconstruction

In the midst of chaos, there is also opportunity.

—Sun Tzu, *The Art of War*

Due to the subatomic nature of the particles considered in this work, one can only directly measure their interactions with detector elements (e.g. energy deposits) as opposed to their direct properties (e.g. their mass). Connecting the detector level observables back down to the fundamental particles of interest is an important step in interpreting the data in the context of the underlying theory. To this end, we identify certain patterns of observables in the detector with the corresponding type of particle we expect to produce them, forming candidate particles. For example, we form candidate electrons (further detailed below in Section 4.3.3) by specific criteria on localized energy deposits within the ECal as well as geometrically matched tracks pointing towards the energy deposits. This procedure is known as object reconstruction and we will frequently refer to detector level signatures identified by this procedure as physics objects.

Relating the predictions of a theoretical model to detector level observables is the inverse problem; we wish to accurately simulate the detector response for events predicted by a certain theory. This is the role of detector simulation. The task of modeling predictions from a specific theory, such as the SM detailed in Chapter 2.2 or a certain BSM theory of Chapter 2.3, is an industry in itself called event generation.

This chapter will discuss these various topics in relating experimental data to theoretical predictions. Section 4.1 details the task of producing representative events from a theory, Section 4.2 discusses the process of detector simulation, and Section 4.3 defines the various criteria for reconstructed particle candidates.

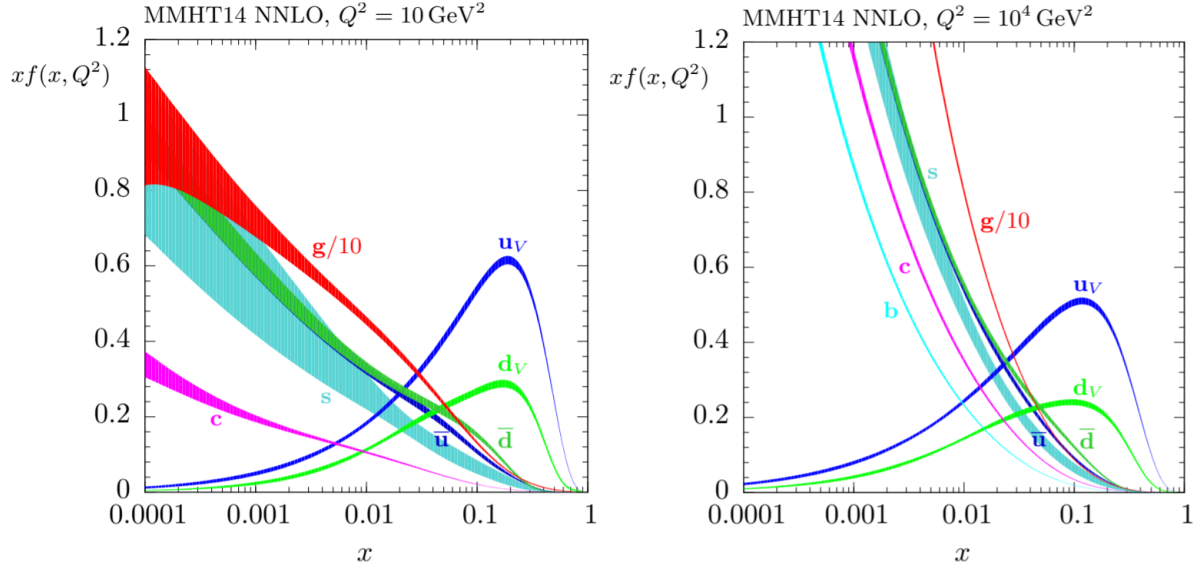


Figure 4.1: Parton distribution functions from the MMHT14 NNLO PDF set for protons at $Q^2 = 10 \text{ GeV}^2$ (left) and $Q^2 = 10^3 \text{ GeV}^2$ (right).

4.1 Event Generation

Interpretation of the results of high-energy physics experiments is heavily reliant on proper modeling of the theoretical predictions of specific processes. In high-energy collider experiments we use tools known as “event generators” to provide this modeling by numerically calculating cross-section and distributions.

One difficulty that needs to be addressed by event generators is the difference in scales and dynamics of the problem over the evolution of an event. In hadron colliders, the initial particles are bound states of several partons undergoing non-perturbative QCD interactions on the timescale of $\hbar/m_p$ [119]. As the incoming partons collide at momentum transfers Q^2 above this value, they interact as effectively isolated particles with no knowledge that they are bound. In addition, the values of the strong coupling constant at $\alpha_s(Q)$ can be in the regime where the interactions and decays can be calculated perturbatively. As the system continues to evolve and individual particles reduce in energy, the non-perturbative effects of QCD will eventually arise and force all the final state particles into colour singlets. This issue of different relevant physics over the scale of the problem is addressed by a factorization ansatz in which these distinct dynamical regions are separated.

The first step of the factorization ansatz is the separation of the matrix element calculation from the initial bound hadrons. As described above, the colliding partons can be treated as effectively isolated so the cross-section for pp collisions to a state X can be factorized as [120]

$$\sigma_{pp \rightarrow X} = \sum_{a,b} \int dx_a dx_b f_a(x_a, \mu_F) f_b(x_b, \mu_F) \sigma_{ab \rightarrow X}(\mu_F, \mu_R) \quad (4.1)$$

where $\sigma_{ab \rightarrow X}$ represents the parton level cross-section and f_a the parton distribution functions (PDF) for a momentum exchange of μ_F (the factorization scale). The PDF represents the probability that the parton a will carry a momentum fraction x_a of the proton. If the factorization assumption of PDFs is valid, they can be measured in one experiment and safely applied to others. Values for the PDFs are compiled by several dedicated collaborations such as NNPD [121, 122], MMHT [123], and CTEQ [124, 125]. An example of some PDF distributions from the MMHT collaboration can be seen in Figure 4.1.

Expanding the matrix element of Equation 4.1 leads to the expression [120]

$$\sigma_{pp \rightarrow X} = \sum_{a,b} \int dx_a dx_b f_a(x_a, \mu_F) f_b(x_b, \mu_F) \frac{1}{2x_a x_b s} \int d\Phi |\mathcal{M}_{ab \rightarrow X}|^2(\Phi, \mu_F, \mu_R) \quad (4.2)$$

where $M_{ab \rightarrow X}(\Phi, \mu_F, \mu_R)$ represents the parton level matrix element and $d\Phi$ is the integration over the phase space of the outgoing particles. An additional scale μ_R is introduced and represents the scale of α_S used in the perturbative expansion. Matrix elements can be calculated perturbatively using Feynman rules, but the final calculation requires integration over the phase space, i.e. all the incoming and outgoing 4-momenta. To calculate this multi-dimensional integral numerically, Monte-Carlo techniques are used, which involves the random sampling of the phase space. Each sampled point which is accepted by the procedure can be interpreted as a single possible event with specific 4-momenta of each particle. The matrix element calculation alongside the PDF allows a calculation of the hard-scatter process.

The outgoing final states may include color-charged states, which undergo gluon radiation or splitting. This is an additional factorization step known as showering which is assumed to be independent on the origin of the colour-charged particle. The differential probability for a color-charged particle a to split to two others, b and c , with z being the momentum fraction of b , and an off-shell propagator with virtuality μ^2 follows the DGLAP equation [126–128]

$$dP_a(z, \mu^2) = \frac{d\mu^2}{\mu^2} \frac{\alpha_s}{2\pi} P_{a \rightarrow bc}(z) dz \quad (4.3)$$

where $P_{a \rightarrow bc}(z)$ is the splitting function for the specific process which is only z dependent. This function is divergent both in the collinear limit, $\mu \rightarrow 0$ and soft limit, $z \rightarrow 0$ (implicit in $P_{a \rightarrow bc}$). These divergences are unphysical as a detector can only resolve emissions at some finite energy and spatial scale. A finite cut-off to the integral range, such as on the relative transverse momenta between the two decays products Q_0 , can thus be imposed to remove the divergences and provide physically meaningful results.

Equation 4.3 is not enough to produce an exact shower algorithm as its product represents all inclusive emissions, and not a specific exclusive multi-emission event. An ordering variable is usually then introduced, such as the virtuality Q^2 (typically set to μ^2). The Sudakov Form Factor then represents the probability of no additional resolvable emissions above Q^2 , but below

a max value Q_{max}^2 [129]

$$\Delta_a(Q_{max}^2, Q^2) = \exp \left(- \int_{Q^2}^{Q_{max}^2} \frac{d\mu^2}{\mu^2} \frac{\alpha_s}{2\pi} \int_{\frac{Q^2}{\mu^2}}^{1-\frac{Q^2}{\mu^2}} dz P_{a \rightarrow bc}(z) \right) \quad (4.4)$$

The term $\Delta_a(Q^2, Q_0^2)$ then represents the probability of no resolvable emission above Q^2 . The derivative of Equation 4.4 is the probability for the first emission. This probability can be re-evaluated for each decayed parton as long as Q^2 is required to get smaller. This provides a full showering algorithm, but introduces some dependence on showering variable, which is a generator specific choice.

The final factorization step is that once the showered states reach the scale of non-perturbative QCD, they undergo a process known as hadronization which binds them into colorless states. Modeling of hadronization is mainly done with phenomenological models such as the “Cluster Model” or the “String Model” [119]. An underlying assumption for these models is that partons do not hadronize independently, but as collective colour-connected states, which undergo the same effects regardless of origin. The implementation of these hadronization models typically involve non-perturbative parameters which can not be calculated by first principle. Such models are tuned by finding which parameter choices best represents observed data. The generator dependent choice of parameters is sometimes known as a “tune”.

Particles with lifetime $c\tau > 10$ mm are considered stable by the event generators. All other particles are decayed until only stable final state particles are found. This is commonly referred to as the “parton-level truth” event, and without loss of generality in this work we will simplify to “truth-level” event.

As described above, event generators use various Monte-Carlo techniques and so are commonly known in the field as “Monte-Carlo Event Generators” or more simply MC (a convention we will use outside this section). Matrix-element generators such as Madgraph [130] and POWHEG-Box [131, 132] are often showered and hadronized by additional generators like Pythia [133] or HERWIG++ [134]. Specific generators such as EvtGen [135] are sometimes used to improve modeling of heavy flavour decays as well.

4.2 Detector Simulation

To provide a direct one-to-one comparison between MC calculations and real data, MC events are passed through a procedure which simulates the detector response of the particles in the event. This task is broken up into two steps: simulation of the final state particle interactions with the detector, and simulation of the digitization of the detector signals [136].

A common choice for the simulation of particles with the detector elements in the high-energy physics community is the GEANT4 [137] tool. A detailed model of the ATLAS detector is provided to GEANT4 which includes the geometry and material properties of the detector el-

elements and support structures. The trajectory of the particles is calculated with a fourth order Runge-Kutta method which includes the effects of the magnetic field in relevant volumes. For each step of the propagation, if a particle passes through a material volume, GEANT4 simulates the expected interaction (energy depositions and/or decays) for the relevant processes including: ionization, bremsstrahlung, photon conversion, Compton scattering, multiple Coulomb scattering, and hadronic interactions. Some of these are discussed in Section 3.2. GEANT4 also simulates the decay of particles which are considered stable in the event generation step, such as π and K mesons.

The output of the simulation step are hits in the detector elements containing space, time, and energy deposition information. Sub-detector specific simulations are applied to convert these values into voltage/current measurements and subsequently into detector digits when these pass hardware thresholds [136]. At this point, the event simulation is in a format directly comparable to the data and can undergo the same object reconstruction detailed in Section 4.3.

The simulation of other interactions beside the main hard-scatter event, such as in-time and out-of-time pile-up, as well as cavern and beam-halo backgrounds are also included in detector simulation step [136]. For these processes, dedicated samples are simulated and added to hard-scatter simulation before the digitization step. The main contribution is from pile-up, which is simulated by the addition of hits from inelastic pp collisions. The number of additional inelastic pp collisions follows a Poisson distribution with mean equal to the $\langle\mu\rangle$ to be simulated.

4.3 Object Reconstruction

The reconstruction of analysis level physics objects relates the detector level observables, such as tracking layer hits and energy deposits, to the underlying physics. Specific algorithms are used to reconstruct physics objects for different particle types, which are detailed in the rest of this section. These objects need to be calibrated to match the properties of the expected particle which caused them, and similarly MC predictions adjusted to match data. Different levels of object definition quality, which usually trade-off lower reconstruction efficiency for lower misidentification rate, called working points, are defined in some cases.

We will make use of a shorthand notation where e , μ , and ν represent an electron, muon, and neutrino respectively and ℓ represents either an electron or muon. τ -leptons are not stable particles and require dedicated reconstruction techniques not detailed in this work. Small- R jets and large- R jets will be identified by j and J which are hadron flavor agnostic, and b will be used for b -tagged small- R jets. Through-out this work, kinematic quantities will be represented using a function notation. For example $p_T(\ell)$ will represent the measured p_T of the lepton. When multiple objects are present, a subscript will be used to order the objects in p_T . The notation j_2 would then represent the sub-leading jet.

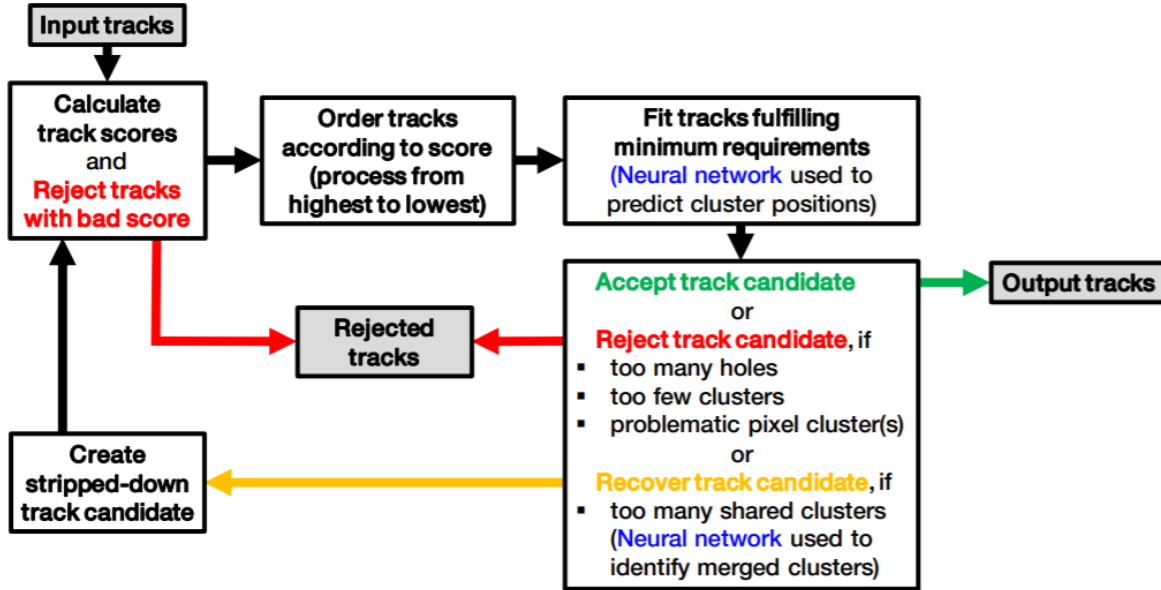


Figure 4.2: Flow chart of the ATLAS track reconstruction procedure [138].

4.3.1 Tracks

As charged particles pass through the ID they leave energy deposits, known as “hits”, in each of the sensitive tracking layers detailed in Section 3.2.1. These hits are combined to form tracks which can be used to reconstruct the charged particle kinematics. This section will detail the algorithm which connects ID hits into tracks, while extrapolation of tracking in the MS system will be discussed in Section 4.3.4.

The main difficulty in track reconstruction in the ATLAS environment is the large charge particle multiplicity, which is a computationally expensive process to untangle. The large quantity of hits can be incorrectly combined to form misidentified tracks, which are not representative of any real charged particle in the event. To reduce these ambiguities, ATLAS uses a tracking algorithm which continuously refines the candidate track collection until only high-quality tracks remain [138]. Figure 4.2 illustrates a flow-chart of the ATLAS tracking procedure.

The first step of the tracking, a clustering algorithm uses a connected component analysis [139] to group SCT and pixel hits with common edges or vertices in a detector layer into a single cluster¹. This is to include the effect of charged particles depositing their energy in several nearby pixels. Groups of three clusters, with requirements on the approximate candidate momenta and impact parameters, are then used to seed a track building algorithm. Track candidates are built from these seeds by applying a combinatorial Kalman filter [140] to match clusters in other tracking layers within a window of the seed. At this point, there can be multiple track candidates per initial seed and track candidates can share clusters. To resolve the possible ambiguities, track candidates are ranked according to a track score, which accounts for

¹In this context two SCT layers are needed to define a specific point and are considered as one layer.

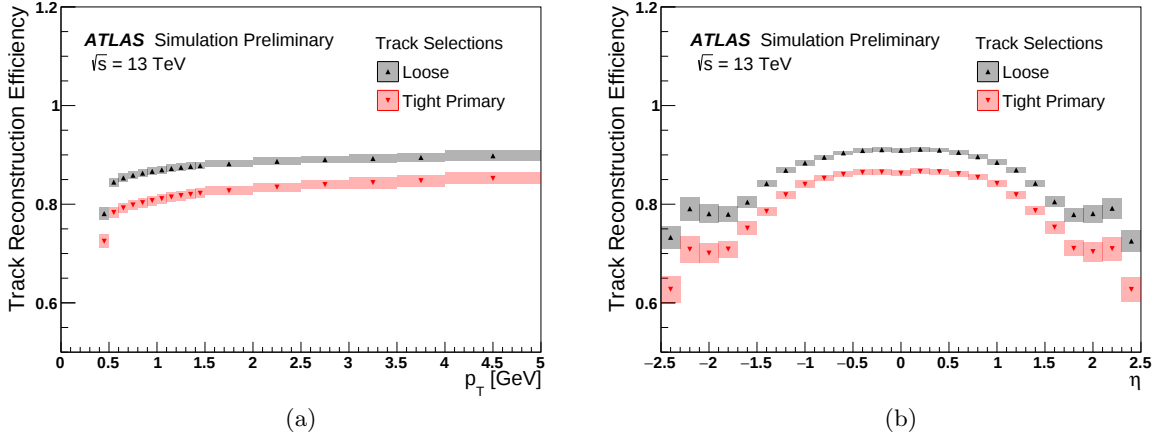


Figure 4.3: Track reconstruction efficiency for the default track definition (black), here called “Loose”, and the tight track criteria (red) as a function of a) p_T and b) η of the track [144].

the hit multiplicity of the track, χ^2 of the fit, and track momenta. Track candidates are then compared to those higher in the ranking, and if they have a shared cluster as identified by a neural-net algorithm [141], the cluster is removed from the lower ranked track which is then re-scored. A final set of criteria on the track hit content, momentum and impact parameter is imposed to define valid track candidates. These additional requirements and those of the track score calculations are detailed in Ref. [138]. The track d_0 impact parameter is defined as the closest approach of the track to the primary vertex (defined below), and z_0 is the longitudinal distance of this point to the primary vertex.

The track candidates are extrapolated to the TRT and combined with those measurements if they pass a track score similar to the ambiguity solving stage [142]. After the tracking is completed an additional outside-in tracking is done with the unused hits, but seeded by TRT tracks via a Hough transform [143]. The outside-in tracking is used to reconstruct tracks which have may have originated from decays away from the IP.

Additional requirements on the tracks can be placed to form “tight” tracks to suppress fake or poorly measured tracks. The full list of additional requirements is detailed in Ref. [144], which amount to a more stringent requirement on the hit multiplicities of the tracks. As shown in Figure 4.3 the efficiency for the usual and tight track reconstruction is approximately 77% and 72% respectively for 500 MeV tracks, increasing to a plateau value above $p_T > 5$ GeV of approximately 90% and 85%.

Tracks are also matched together with an annealing algorithm to reconstruct candidate vertices, representing particle decays or interactions [145, 146]. The algorithm implements an adaptive vertex finding method [147] to sequentially find tracks consistent with a common origin. Iteratively, the vertex position is calculated by a χ^2 minimization of all tracks, followed by a reweighting of the tracks based on their closest approach to the vertex. Once the annealing

is finished, tracks greater than seven standard deviations from the vertex are removed, and the procedure continues with the removed tracks until no more vertices can be formed. Vertices are required to have at least two tracks associated to them, and the one with the highest $\sum p_T^2$ is identified as the primary vertex (PV).

4.3.2 Jets

Colour-charged particles undergo a showering process as discussed in Section 4.1. The observed final state of a single energetic colored particle then includes a multiplicity of particles which we wish to relate back to the parent particle which initiated the shower. To reconstruct these multi-particle showers, elements of detector observables, such as energy deposits, are clustered into objects we call “jets” using a specific “jet algorithm.”

We will consider jets clustered by the anti- k_t algorithm [113]. The anti- k_t algorithm is both collinear and infrared safe, which indicates that it is insensitive to the collinear and soft divergences of QCD showering (see Section 4.1). The 4-momentum of the anti- k_t jet is then a good representation of the original colour-charged particle which produced the shower. In addition, the anti- k_t algorithm has the property that the boundary and area of the jet is insensitive to soft radiation, and hence less affected by hadronization and pile-up effects. The only parameter for the anti- k_t algorithm is a radius value R , which approximately governs the size of the jets. For hard particles separated by more than R , the anti- k_t algorithm will group all soft particles within a cone of this size around the hard particles. For hard-particles closer than R , the jet shapes become more complicated.

Various detector observables can be clustered to form jets. We will use several different jet definitions, which we will call jet collections, throughout this work and will be explicit when needed. A default collection of small- R jets using calorimeter clusters will be defined for reconstruction of individual quark and gluons. A subset of the small- R jets collection will be those identified as the those containing b -hadrons, which we will called b -tagged jets. For identification of high- p_T decays of $V \rightarrow q\bar{q}$, a separate collection of large- R calorimeter jets will be used due to finite resolution of small- R jets. To assist in identifying b -hadrons with the large- R jets, a new collection using solely tracking information, the track-jets, will also be used.

Small- R Jets

When referring to jets, without any additional modifier, we will in particular refer to $R = 0.4$ anti- k_t jets clustered on topologically clustered cells² calibrated to the EM scale (EMTopo) [114]. This jet collection is the most useful for general physics cases and can be used to measure quarks and gluons from the hard-scatter process. We will refer to these as small- R jets when distinction is needed.

²Note the different usage of the term cluster. The jet clustering algorithm is applied to objects which are clusters of cells.

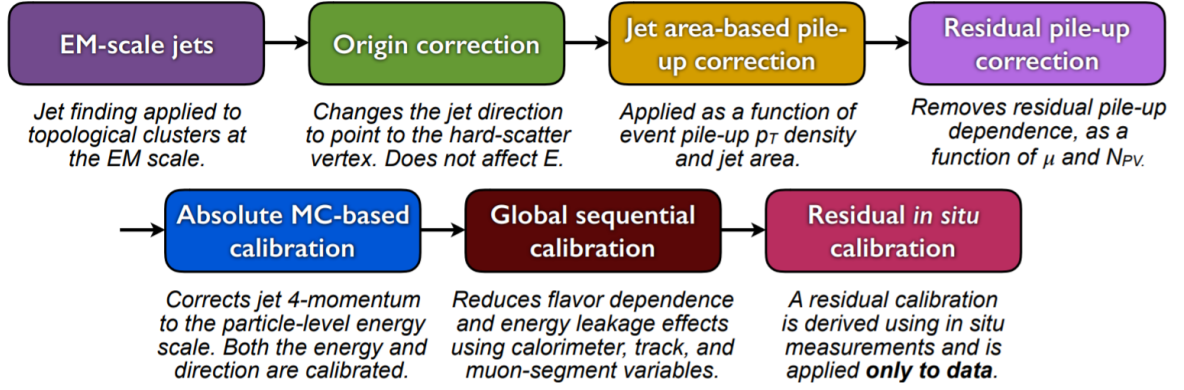


Figure 4.4: Flow chart of the jet energy scale calibration [148].

The topological-clustering algorithm is a spatial matching of calorimeter cells, seeded and expanded based on the local significance of the cell energy with respect to noise. For each cell the local significance ζ_{cell}^{EM} is defined as the ratio of the energy of the cell E_{cell}^{EM} and the standard deviation of the cell noise $\sigma_{cell,noise}^{EM}$. The cell noise includes both the effects of electronic noise and pile-up, which are estimated as Gaussian fluctuations in data and MC respectively. Seed cells, defined as having $|\zeta_{cell}^{EM}| > 4$, initiate the clustering process and all neighboring cells in the detector layer or in adjacent layers with η or ϕ overlap, are grouped into the cluster. If any of these cells have $|\zeta_{cell}^{EM}| > 2$, the process continues, and the following neighboring cells are also added to the cluster. If cells are shared by multiple clusters, the clusters are merged. The clusters formed in this way can become too large, including the contributions from several nearby particles, so an additional cluster splitting step is done. Cells which are local maximum and have $E_{cell}^{EM} > 500$ MeV are used to split the cells evenly between the maxima in all spatial directions. The final 4-momenta of the cluster is then formed using a energy weighted barycenter calculation for η and ϕ , the total energy of all cells E_{clus}^{EM} , and the assumption of zero mass. Only clusters with positive E_{clus}^{EM} are used in the jet algorithms, but $|\zeta_{cell}^{EM}|$ is used during the clustering to average out the effect of in-time and out-of-time pile-up.

The EMTopo cluster and cell energies above are measured and calibrated at the EM scale and do not include the effects of signal loss in the non-compensating calorimeters for hadronic signals. To correct for these effects, the jet energy scale (JES) for EMTopo jets is calculated and used to calibrate the jet energy back to the true energy scale [148, 149]. The JES calculation is done over several steps shown in Figure 4.4.

The first step is an origin correction which recalculates the cluster η and ϕ to point to the hard-scatter vertex as opposed to the ATLAS geometric center. Next, a per-event pile-up subtraction is applied where the pile-up p_T density of the event, ρ , is evaluated and multiplied by the area of the jet A [150] to determine the pile-up contribution to the jet p_T . Only the low occupancy $|\eta| < 2$ clusters are used for the pile-up subtraction step, which results in a non-optimal subtraction in the forward region. To compensate for this effect, a residual pile-up

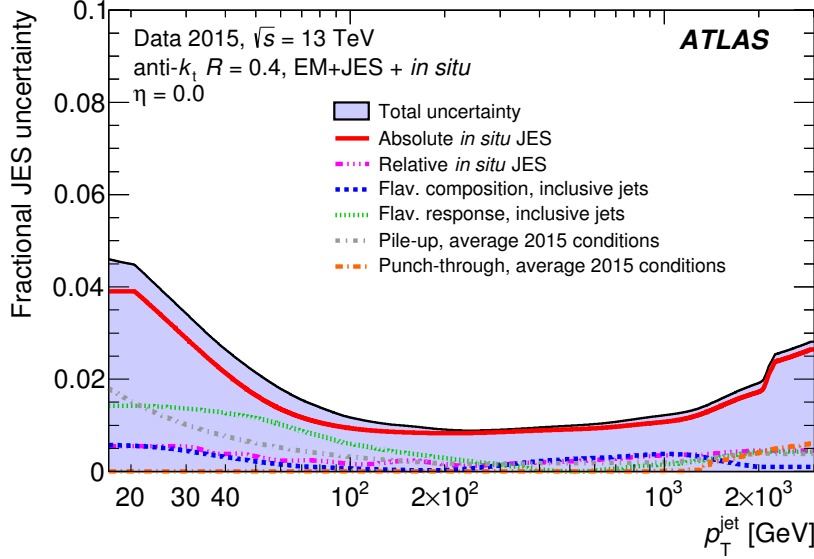


Figure 4.5: The uncertainty on the small- R JES as a function of jet p_T . The blue filled region shows the total uncertainty and the curves show individual contributions. The dominant contribution originates from the *in-situ* calibrations [148].

correction dependent on the number of primary vertices N_{PV} and pile-up μ , is calculated as a function of jet p_T and η . The absolute JES is calculated by matching truth-level jets to reconstructed jets within $\Delta R < 0.3$. The jet energy response $R_E = \frac{E^{EM}(jet)}{E^{truth}(jet)}$ is calculated as the ratio of reconstruction jet energy to the truth jet energy. The JES is then calculated in bins of jet energy and η , and applied to reconstructed jet energies so that $R_E = 1$. The JES evaluated here is inclusive of all jets and can show residual effects based on the structure of the jet (example a quark or gluon initiated jet). To correct for these effects, a global sequential calibration is applied, where a list of observables sensitive to these effects (provided in Ref. [148]) are sequentially calibrated such that the absolute JES is unchanged, while enforcing a uniform R_E in the observable as well.

The previous JES calibration steps are all derived from MC simulations. To further calibrate the effects of detector modeling and simulation, a final data-driven *in-situ* calibration is derived to be applied on data. Due to momentum conservation in the transverse direction, the jet response can be accurately measured in events where a jet is found recoiling against better calibrated objects. The first of the *in-situ* measurements is an η -intercalibration, where the response of forward $0.8 < |\eta| < 4.5$ jets are balanced against central jets $|\eta| < 0.8$, which allows the correction of forward jet response to match those of central jets. Next the average jet response $\langle R_E \rangle$ is measured in events with jets recoiling off a well-measured photon or Z , providing calibrations up to jet $p_T \approx 1$ TeV. Lastly events with high- p_T jets ($300 < p_T < 2000$ GeV) recoiling against several well measured low- p_T jets are used to calibrate the high p_T jet region. The relative uncertainty on the p_T from the JES calibration at $|\eta| \approx 0$ can be see in Figure 4.5.

The JES calibration involves a total of 80 systematic uncertainties propagated from the individual steps detailed above, 67 of which originate from the *in-situ* calibrations. To reduce the total number of sources of uncertainty to be evaluated for an individual analysis a nuisance parameter reduction scheme is done by applying an eigen-decomposition of the correlation matrix of the *in-situ* uncertainties. A N -parameter reduction scheme is derived by retaining only the N eigenvectors with the largest eigenvalues, with a single additional residual uncertainty constructed to approximate the remaining systematics. This scheme approximately conserves the total uncertainty and the correlation effects. A $N = 21$ parameter scheme will be used in Chapters 5 - 7.

After the JES calibration, the jet energy resolution (JER) is also calibrated [151, 152]. The noise term is evaluated by measuring the energy balance in events one-bunch crossing after a calorimeter-based trigger, and the stochastic+constant term are measured using the same *in-situ* γ/Z -jet and dijet balance techniques as the JES. The total number of systematics from the JER evaluation are reduced using a similar eigen-decomposition as implemented for the JES. The relative uncertainty on the jet p_T from JER calibration for jets with $|\eta| \approx 0.2$ is approximately 2.5% at 20 GeV and less than this value from 70 - 2000 GeV.

To reduce the number of misidentified jets originating from beam-induced-backgrounds, cosmic-ray muons, and calorimeter noise, additional requirements on the energy distribution of the cells within the cluster and the associated tracks are applied to form a “cleaned” set of jets [153]. The jet cleaning has greater than 99.5% selection efficiency for jets originating from hard collision events. In almost all circumstances jets are also required to have $p_T > 20$ GeV in order to be considered well measured.

The jet cleaning removes a significant number of misidentified and poorly measured jets from non-collision backgrounds, but additional pile-up collisions can also produce a large multiplicity of jets which are not relevant to the hard-scatter process. To suppress pile-up jets, a JetVertexTagger (JVT) [154] variable is used for jets within the tracker acceptance of $|\eta| < 2.4$. By utilizing the tracking information, the amount of energy in a jet from pile-up vertices can be estimated by measuring the number of tracks originating from those vertices that are matched to the jet clusters. Two track-based observables (detailed in Ref. [154]), which measure the contribution of a jet’s momentum originating from the PV are provided as inputs to a k-Nearest Neighborhood algorithm [155], with the result designated the JVT score. With this algorithm, the JVT score can be interpreted as the probability of a jet originate from the hard-scatter vertex. For $p_T < 60$ GeV and $|\eta| < 2.4$ a JVT score > 0.59 is required to suppress pile-up jets, which corresponds to an approximately 92 (1)% signal (background) jet efficiency [154].

Unless otherwise specified, when referring to small- R jets, we will refer to cleaned jets with $p_T > 20$ GeV, $|\eta| < 4.5$, and the JVT criteria applied for jets within $|\eta| < 2.5$.

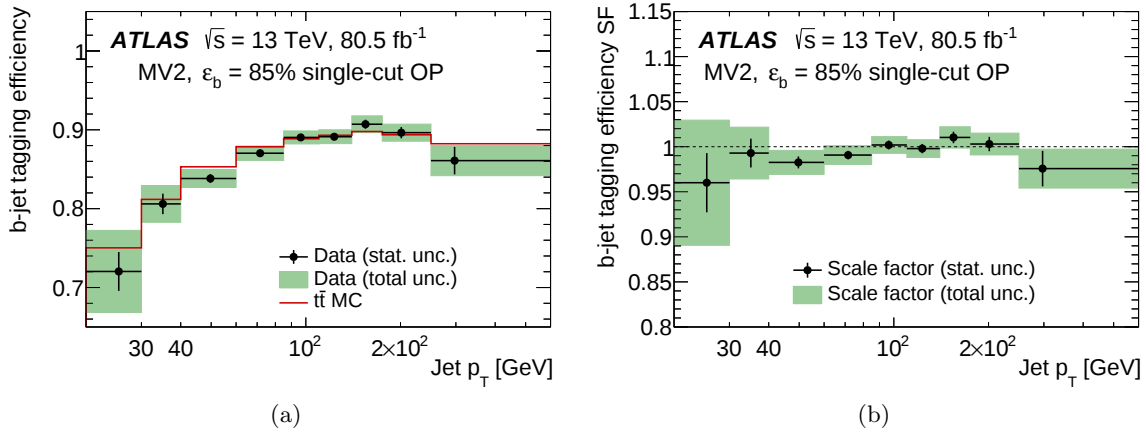


Figure 4.6: a) The b -tag efficiency and b) MC scale-factor for the 85% efficiency working point as a function of jet p_T [157]. The extrapolation to high- p_T regime is not shown. The black points show the data and statistical uncertainty and the green filled regions show the total uncertainty on data. The $t\bar{t}$ MC prediction used to derive the scale-factor can be seen in red in a).

b -tagged Jets

Hadrons containing b -quarks differ from other hadrons as they typically have larger masses, lifetimes, and multiplicities of decay products. A 50 GeV b -hadron travels on average 3 mm before decaying, which results in visible secondary vertices [156]. Identification of b -hadrons is useful for many physics processes such as identifying t -quarks, which decay into a W -boson and a b -quark greater than 95% of the time [88].

The ATLAS experiment has developed several techniques to identify jets containing b -hadrons, known as b -tagged jets, from the light flavour jets (u, d, g -induced jets), and c -jets. Jets with c -hadrons have characteristics intermediate to those of b -jets but and light quark jets. The overall approach is to identify several low-level variables, mainly those related to the displaced decays of b -hadrons, which are then combined into single discriminant with machine learning techniques [156, 157]. The low-level algorithms make use of the kinematics and secondary vertices formed from tight tracks with $p_T > 500$ MeV (see Section 4.3.1) which are within a p_T dependent ΔR window to jets and not consistent with Λ baryons, K mesons, or photon conversions. The low-level taggers either attempt to identify b -jets based on the distribution of impact parameters of the associated tracks, or by directly attempting to reconstruct the secondary vertices of the b -hadron decay.

The low-level tagger information, alongside the jet p_T and $|\eta|$, is then provided to a Boosted Decision Tree algorithm, known as MV2c10, and trained to provide the best separation between, b -jets and other jets. Several working points are defined by the b -jet signal efficiency for a fixed cut on the MV2c10 score. The 85% signal working point, has rejection factor of 25 (2.7) for light-flavour jets (c -jets) [157]. The b -tagging efficiency as a function of jet p_T can be seen in

Figure 4.6a for the 85% working point as estimated in $t\bar{t}$ events. When referring to b -tagged jets in subsequent chapters, we will refer to jets identified with the 85% working point of the MV2c10 tagger.

Scale factors are applied to MC to include differences in the b -tagging efficiency with respect to data. The scale-factors are measured in data for light-flavour, c -jets, and b -jets independently using a $t\bar{t}$ enriched sample [157–159]. The b -tag scale-factor as a function of jet p_T for the 85% working point can be found in Figure 4.6b as estimated in $t\bar{t}$ events.

The b -tagging systematics are those related to uncertainties on the efficiency calculation, which are mainly from generator modeling uncertainties and propagation of experimental uncertainties. The efficiency uncertainties are evaluated separately for light-flavour, c -jets, and b -jets. To reduce the total number of systematics propagated to the analyses, an eigenvalue decomposition of the uncertainties is done using the same techniques as the JES [157]. For most of this work a reduction scheme is used where the top three/four/five eigen-parameters are retained for the b/c /light jets respectively. An additional extrapolation uncertainty is applied in the $p_T > 400$ GeV regimes where the calibrations do not have significant statistics for reliable estimates.

Large- R Jets

W/Z bosons have decay modes to two quarks, with the opening angle between the two quarks following the relation

$$\Delta R \approx \frac{2m(V)}{p_T(V)} \quad (4.5)$$

where $m(V)$ and $p_T(V)$ are the mass and momentum of the boson respectively. As the p_T of the boson gets larger it becomes increasingly difficult to reconstruct the out-going quarks as individual $R = 0.4$ jets. Instead it becomes useful to consider reconstruction of the W/Z boson with a single jet with a larger radius parameter. For bosons with $p_T > 200$ GeV it is expected that most of the hadronic shower can be reconstructed with $R = 1.0$ jets, called large- R jets.

Large- R jets reconstructed using a $R = 1.0$ anti- k_t algorithm use a different approach for the calibration with respect to small- R jets. Individual EMTopo clusters are first calibrated to the hadronic scale using a “local hadronic cell weighting” (LCW) approach, where the possible contribution of hadronic showers is evaluated per cluster and the corresponding cell signals reweighted. The probability that a cluster is from an EM shower, P_{clus}^{EM} , is determined from the cluster depth λ_{clus} (distance from the cluster center of gravity to the front of the calorimeter), and the average cluster density ρ_{clus} (the cluster average of $\rho_{cell} = E_{cell}^{EM}/V_{cell}$, where V_{cell} is the cell volume) The reconstruction level probabilities are calculated from simulated pion events. The cells are then weighted to a new energy scale by:

$$w_{cell}^{cal} = P_{clus}^{EM} \cdot w_{cell}^{EM} + (1 - P_{clus}^{EM}) \cdot w_{cell}^{HAD} \quad (4.6)$$

where w_{cell}^{EM} is assumed to be 1, and w_{cell}^{HAD} is the cell correction as evaluated from the same

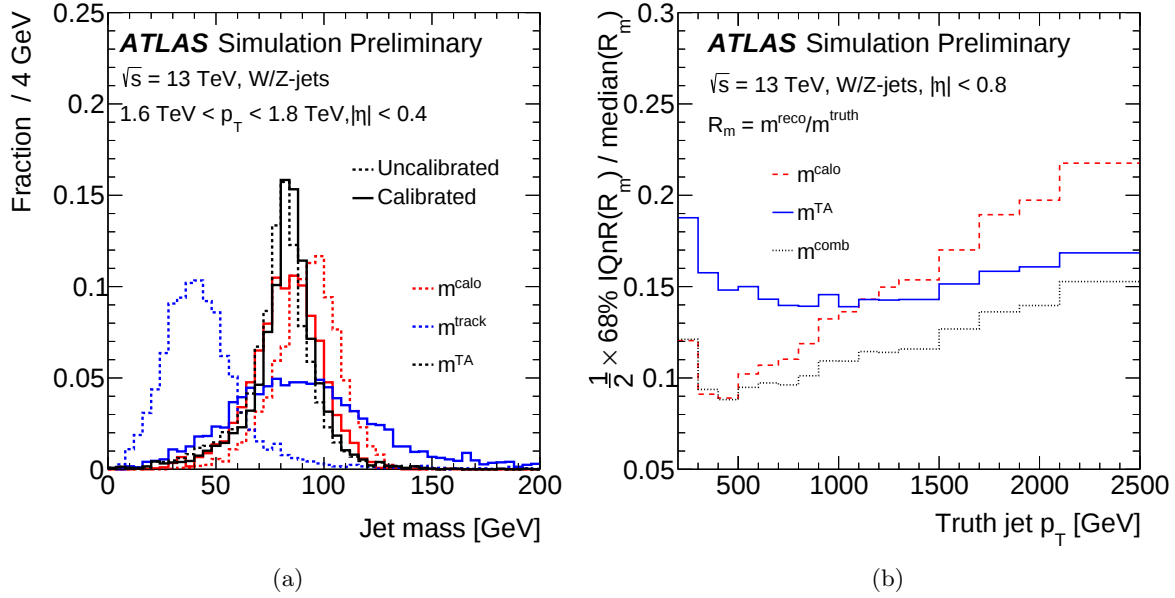


Figure 4.7: a) Jet mass as measured in 13TeV V +jets MC using a calorimeter mass definition (red), track mass definition (blue), and a track assisted mass (black). Both calibrated (solid) and un-calibrated (dashed) values are shown b) The large- R jet mass resolution as a function of jet p_T for the calorimeter mass (red), track assisted mass (blue), and combined mass (black) definition. The mass resolution is quantified as half the interquantile ratio divided by the median jet mass, which is an outlier invariant measure of the resolution. The combined mass shows the best resolution over the whole p_T range [163].

set of MC. The $w_{\text{cell}}^{\text{HAD}}$ calibrations are binned in cell layer, η , ρ_{cell} , and $E_{\text{clus}}^{\text{EM}}$. The effects of out-of-cluster corrections and dead material corrections are applied as additional weights over the whole cluster. These new locally-calibrated topologically clusters (LCTopo) are used as inputs for the $R = 1.0$ jets.

Due to the larger radius parameter, $R = 1.0$ jets will typically have larger areas and be more susceptible to pile-up. To remove the pile-up contribution, a jet trimming [160] procedure is applied. The jet trimming is a specific type of “grooming algorithm” which is found to be the most effective in the context of W -identification [161]. In the trimming procedure, the clusters associated to the jet are re-clustered into sub-jets with a k_t algorithm with a radius parameter $R_{\text{sub}} < 1.0$. Sub-jets with momentum less than some fraction f_{cut} of the initial jet p_T are removed. The choice of $R_{\text{sub}} = 0.2$ and $f_{\text{cut}} = 0.05$ are found to be optimal for ATLAS studies [161].

The jet trimming acts as a substitute for the pile-up based subtraction detailed in the small- R jet JES calibration. The large- R jet JES and JER are calculated in the similar fashion as the small- R jets detailed above [162]. In particular, the same set of *in-situ* Z/γ and dijet balance techniques are used.

The mass of the large- R jet is also a useful quantity so the jet mass scale (JMS) and jet

mass resolution (JMR) are also measured after the JES calibration. A different set of *in-situ* techniques are used for calibration of the JMS. The first of these is the R_{trk} method where the ratio of calorimeter-jet to track-jet p_T , $R_{trk} = \left\langle \frac{p_T^{calo}}{p_T^{trk}} \right\rangle$, is measured in data and dijet MC and used as a scale-factor. The second method implements a forward folding technique on samples of semi-leptonic $t\bar{t}$ events through fits to the hadronic top-quark and W boson mass and widths in MC and data. The relative uncertainty on the JES and JMS varies from approximately 2% at jet p_T of 200 GeV, up to 6 - 8% at 3 TeV.

The JMR is found to be improved by utilizing a track assisted calculation of the jet mass in addition to the usual calorimeter based quantity. Track jets are clustered from tight tracks with $p_T > 500$ MeV using the same clustering+trimming algorithm and ghost-associated [150] to the calorimeter jets. The track assisted mass is then calculated as [163]

$$m_{TA} = m_{trk} \frac{p_T^{calo}}{p_T^{trk}} \quad (4.7)$$

where p_T^{calo} and p_T^{trk} are the matched calorimeter and track jet p_T . As shown in Figure 4.7, the m_{TA} definition is found to have better resolution for high- p_T jets, which is when the calorimeter resolution becomes too coarse to resolve the individual particle contributions. At low- p_T the track assisted mass begins to perform worse as the calorimeter provides better energy resolution than the tracking system. To provide the best overall resolution, the two mass terms are weighted to form a combined mass definition [163]

$$m_{comb} = w_{calo} \cdot m_{calo} + w_{TA} \cdot m_{TA} \quad w_{calo} = \frac{\sigma_{calo}^{-2}}{\sigma_{calo}^{-2} + \sigma_{TA}^{-2}}, \quad w_{TA} = \frac{\sigma_{TA}^{-2}}{\sigma_{calo}^{-2} + \sigma_{TA}^{-2}} \quad (4.8)$$

where σ_{calo} and σ_{TA} are the calorimeter and track jet mass resolutions estimated from the 68% interquantile range of the jet mass response. The m_{comb} definition provides the best overall mass resolution, as seen in Figure 4.7.

At the analysis level, when referring to large- R jets we will refer to jets clustered with $R = 1.0$ anti- k_t and trimmed as above. Only jets with $p_T > 200$ GeV and $|\eta| < 2.5$ will be considered. The combined mass in Equation 4.8 will be used in all cases.

Large- R jets from W/Z decays have the distinct feature in their radiation pattern where most of the energy deposits are localized in two sub-regions associated with the two outgoing quarks. Features such as this are known as jet-substructure, which are useful for discriminating against background jets which lack them. The amount of substructure in the jet can be measured by the two and three-point energy-correlation-functions within the jet defined as

$$ECF_2^\beta = \frac{1}{p_T^2(J)} \sum_{1 \leq i \leq j \leq n} p_{Ti} p_{Tj} (\Delta R_{ij})^\beta \quad (4.9)$$

$$ECF_3^\beta = \frac{1}{p_T^3(J)} \sum_{1 \leq i \leq j \leq k \leq n} p_{Ti} p_{Tj} p_{Tk} (\Delta R_{ij} \Delta R_{ik} \Delta R_{jk})^\beta \quad (4.10)$$

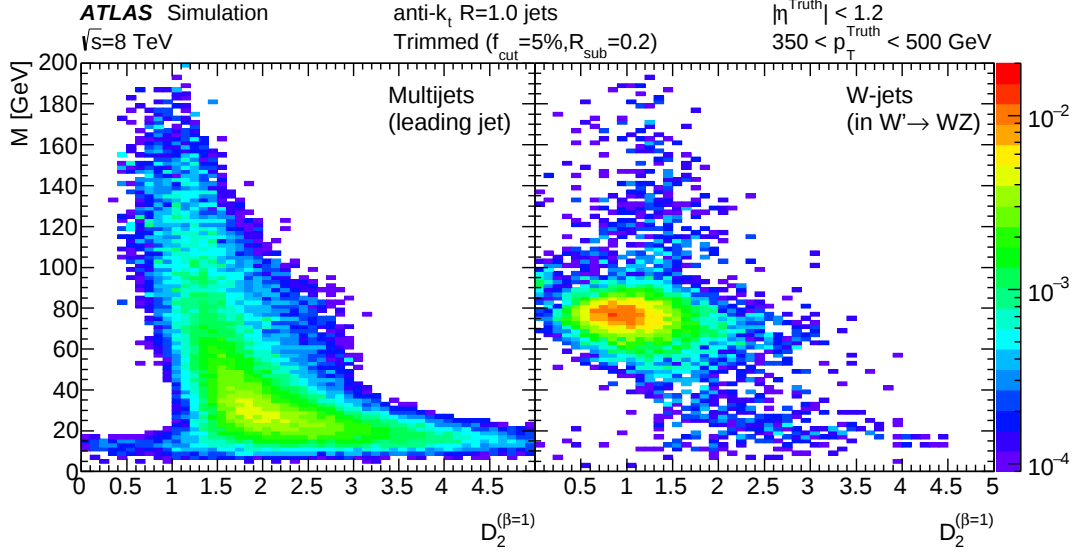


Figure 4.8: Normalized distribution of trimmed large- R jets as a function of the invariant mass and $D_2^{\beta=1}$ as measured from 8 TeV MC from a) multi-jets and b) hadronically decaying bosons [164]. The z -axis shows the log of the binned probability for each of the sources. Only jets with $|\eta| < 1.2$ and $350 < p_T < 500$ GeV are selected.

where the sums are over all the clusters associated to the large- R jet. Here p_{Ti} represents the cluster energy, and ΔR_{ij} the relative opening angle ΔR between cluster i and j . The β is an arbitrary scaling value. The power-law scaled ratio of these correlation functions defined as

$$D_2^{\beta=1} = \frac{ECF_3^{\beta=1}}{\left(ECF_2^{\beta=1}\right)^3} \quad (4.11)$$

was predicted theoretically to provide strong discriminating power for W/Z jets against jets of different origin [165, 166]. In comparison to other substructure metrics, the $D_2^{\beta=1}$ variable was found in ATLAS data and simulation to offer the best discriminating power for LCTopo large- R jets to identify W/Z jets [161, 164, 167]. Figure 4.8 shows a plot of the $D_2^{\beta=1}$ and invariant mass distribution of large- R jets originating from multi-jets and W/Z boson decays. The distribution of these variables will be used in Section 5.2.3 to design a V -tagger to identify large- R jets consistent with a W/Z boson.

Track-jets

Utilizing b -tagging with large- R jets faces several issues as most of the b -tagging algorithms exploit very specific features of single b -hadron decays. For large- R jets, several b -hadrons can be expected, for example in $Z \rightarrow b\bar{b}$ decays. In addition, the trimming procedure, which attempts to remove pile-up contributions, may remove necessary information for the b -tagging algorithms.

A technique to avoid these issues is to apply jet-clustering on tracks and apply b -tagging techniques directly on this jet-collection. The “track-jets” can then be ghost-associated [150] to the large- R jets and the b -tagging information of the track-jets can be used to infer information on possible b -hadron constituents with the large- R jet.

Since track-jets are a separate jet-collection with respect to calorimeter jets, they can be optimized and calibrated independently for optimum b -tagging performance. Parameters can be chosen to utilize the finer granularity tracking system in comparison to the calorimeter. It is found that track-jets clustered with the anti- k_t algorithm with $R = 0.2$ provide the best background rejection for multiple b -jet boosted topologies [168].

Only tight tracks with $p_T > 400$ MeV from the PV are utilized in the jet algorithm. Track-jets are required to have $p_T > 25$ GeV and $|\eta| < 2.5$ to be considered. The MV2c10 algorithm training, calibration, and data/MC scale-factors are calculated also for track-jet usage.

4.3.3 Electrons

As charged stable particles, electrons will leave both tracks in the ID as well as clusters in the ECal. Candidate electrons are then identified as calorimeter deposits and tracks which are both consistent with an electron and can be matched. The matching is important for separation of electron candidates with photons, which have similar EM showering behavior but no associated tracks³.

ATLAS uses a two-step electron identification procedure where seed clusters and candidate tracks are first identified, then a second step which combines the two and refines them into a final electron definition [169]. To identify seed clusters, the ECal is segmented into intervals of $\Delta\eta \times \Delta\phi = 0.025 \times 0.025$, and summed across all the layers to form calorimeter towers [169, 170]. A sliding window algorithm is then applied on the calorimeter towers to identify local maxima [171]. Windows of size 3×5 in $\Delta\eta \times \Delta\phi$ scan over the calorimeter and any local maxima above 2.5 GeV are identified as seed clusters. If two maxima overlap in a 5×9 window, the cluster with largest central tower E_T is kept if both energies are within 10% of each other, or the one with largest cluster E_T is kept otherwise.

Electrons can suffer from significant bremsstrahlung as they transverse the ID, with the radiated photons possibly converting to electron-positron pairs. Multiple track candidates can then be matched to one seed/EMTopo cluster. The first pass of the cluster-track matching scheme tries to identify loose tracks (See Section 4.3.1) which can be extrapolated in a loosely defined region of interest around a seed cluster. If no such track is found, a second pass is attempted where the track fit is redone allowing up to 30% energy loss per-layer crossed. Track candidates are then fit with a global χ^2 algorithm [172]. To further account for the worsened resolution due to radiative losses, the parameters of track candidates are recalculated using a

³Photons can also be “converted” into electron-positron pairs by interactions with detector elements before the calorimeter. Such photons can still be partially identified but will not be discussed in the context of this work.

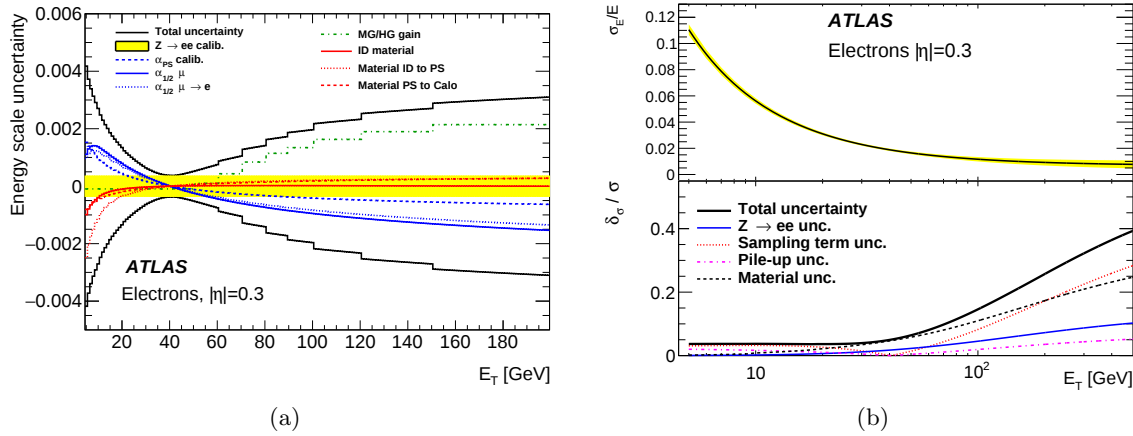


Figure 4.9: a) The uncertainty on the electron scale as function of E_T . The total uncertainty is shown by the black symmetrized envelope and the relative contributions from different individual sources are shown separately. b) The uncertainty on electron energy resolution as a function of E_T at $|\eta| \approx 0.3$. The top plot shows the total uncertainty and the bottom plot shows the relative size of individual contributions [174].

Gaussian-Sum Filter (GSF) [173] to include the effects of experimental noise and non-linear tracking effects.

The primary track associated to the seed cluster candidate is chosen based on the number of pixel hits and the angular matching of the track-cluster pair as detailed in Ref. [169]. After additional requirements on the primary track hit content, the calorimeter window around the candidate electron cluster is widened to 3×7 in $|\eta| < 1.37$ or 5×5 in $1.52 < |\eta| < 2.47$, forming the final candidate electron.

The calibration of the reconstructed electron energy scale and resolution is done with a multivariate regression using a Boosted Decision Tree (BDT) algorithm trained on $Z \rightarrow e\bar{e}$ events [174, 175]. Scale-factors are derived during the calibrations to compensate for differences in the distributions predicted by MC and those of data. Figure 4.9 shows the uncertainty on the electron energy scale and resolution. The uncertainty from the calibration is less than 0.2% for 45 GeV electrons and $|\eta| \approx 0.3$.

The candidate electron definition above provides an efficiency greater than 97% for candidates with $p_T > 15$ GeV [169], but a large number of misidentified or non-prompt electrons can also pass this definition. To suppress such backgrounds, additional identification quality and isolation requirements are typically placed. As detailed in Ref. [169], three identification quality working points “loose”, “medium”, and “tight” are defined using a likelihood based discriminant to separate signal electrons from photon and jet backgrounds. These provide respectively signal efficiencies of 93%, 88% and 80% at $E_T = 40$ GeV and plateau to 96%, 94% and 90% at higher energies. These working points are expected to have background efficiencies of 0.8%, 0.5% and 0.3% respectively at 20 GeV, decreasing to 0.3%, 0.015%, and 0.01% at 70 GeV [176].

Isolation requirements quantify the amount of detector activity around the candidate elec-

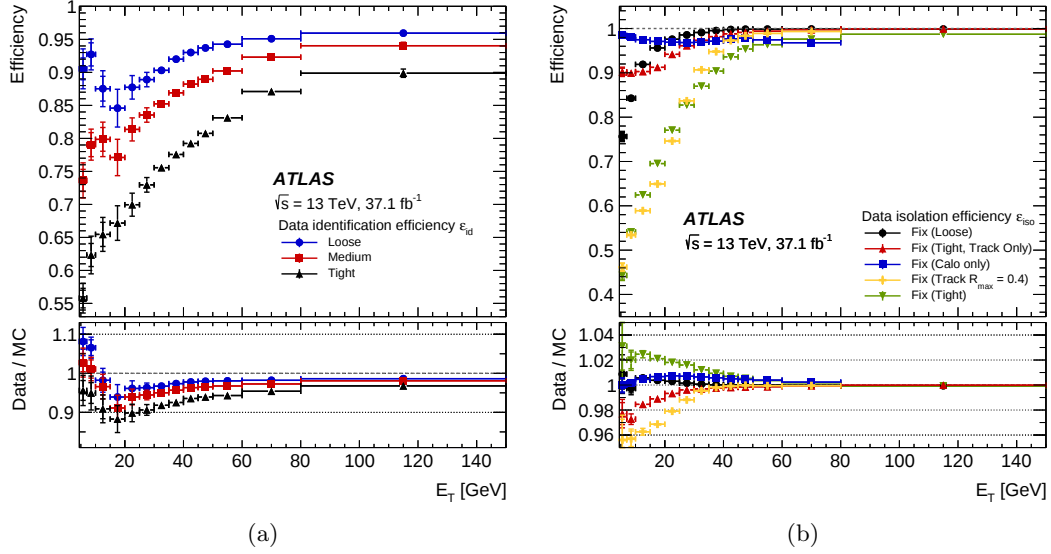


Figure 4.10: The efficiency in data of a) loose, medium, and tight electron quality definitions and b) different electron isolation working points in $Z \rightarrow e\bar{e}$ events. The ratio of the efficiencies as estimated between data and MC are shown in the bottom sub-plots [174].

tron. Candidate electrons which are more isolated are typically better measured and more likely to originate from real electrons. Isolation criteria can be defined using either calorimeter or tracking information [176]. The calorimeter isolation E_T^{iso} is defined as the sum of all topo-clustered energy deposits within $\Delta R = 0.2$, with the core 3×7 cells removed. An additional area-based pile-up subtraction and empirically measured leakage correction are also applied to the calorimeter based isolation calculation [176]. The track-based isolation $p_T^{iso,var}$ is calculated as the sum p_T of tracks with $p_T > 1$ GeV, $|z_0 \sin \theta| < 3$ mm and within the electron p_T dependent window $\Delta R = \min(\frac{10\text{GeV}}{p_T[\text{GeV}]}, 0.2)$. Tracks which are associated to the electron candidate or are within a $\Delta\eta \times \Delta\phi < 0.05 \times 0.1$ window when extrapolated to the cluster layer are also removed. A series of isolation working points is derived in Ref. [176] for varying performance requirements. In this work we will use the “FixedCutTight” isolation definition, which applies cuts on E_T^{iso} and $p_T^{iso,var}$ as well as “LooseTrackOnly” which places E_T and η dependent cuts on $p_T^{iso,var}$ for uniform efficiency. The LooseTrackOnly isolation requirement is 99% efficient for all electrons kinematics, while the FixedCutTight efficiency is 95% efficient for the signal electrons considered in Chapters 5 - 7. Figure 4.10 displays the E_T dependent efficiencies and scale-factors of the identification requirements and different FixedCut isolation requirements.

4.3.4 Muons

Muons are charged particles which typically punch-through the calorimeter since they do not deposit a large amount of their energy into showers. They are then expected to have tracks both in the ID, and in the MS system. The formation of ID tracks are detailed in Section 4.3.1,

while this section will detail how MS tracks are formed and muon candidates defined.

MS tracks are first constructed by forming individual track segments in each of the sensitive layers [177, 178]. In the MDT, a Hough transform is used to form linear track segments from the drift-tubes hits, and the RPC or TGC is used to measure the track segment coordinate in the bending plane. The CSC does a combinatorial search in the two readout planes, with a loose constraint on the segments to be pointing from the IP. Full track candidates are then built by a seeded algorithm which matches segments based on relative positions, hit multiplicity and quality. The track candidates are first seeded by segments in the middle layer, which has a high multiplicity of trigger chambers, followed by the outer two layers. Track candidates are required to have at least two segments except in the transition regions. Finally all the hits forming track candidates are then fit with a global χ^2 fitter [172] to form MS tracks. Hits which provide a large χ^2 can be removed during the fit, and additional hits added if they are consistent with the track.

Final muon candidates are then constructed by combining the MS tracks with information from other detectors. “Combined muons” are constructed from a global refit of the ID and MS tracks, first by back-extrapolating MS tracks, then forward-extrapolating ID tracks. “Segment-Tagged” muons are formed by matching forward-extrapolated ID tracks to single MS track segments. “Calorimeter-tagged” are similar to the segment-tagged but the ID track is matched to calorimeter deposits consistent with a muon. Overlaps are solved by prioritizing Combined Muons, followed by Segment Tagged. Leftover MS tracks which are loosely consistent with the IP are denoted as “Extrapolated” muons. Only certain subsets of these muon definitions are allowed in the muon identification working points discussed below.

Differences in candidate muon momentum scale and resolution in data and MC are calibrated with scale-factors derived from tag-and-probe methods with $Z \rightarrow \mu\mu$ and $J/\psi \rightarrow \mu\mu$ decays [177, 178]. The uncertainties on these scale-factors are 2% for 10 GeV muons and 0.4% for 100 GeV [177].

Similar to electrons, additional identification and isolation requirements are placed on muon candidates to reduce misidentified backgrounds or non-prompt muons. Muon identification is set by three working points, “loose”, “medium”, and “tight”, which apply selections on the candidate muon algorithm p_T , η , hit quantities, fit χ^2 , and differences in ID and MS track p_T detailed in Ref. [177]. The approximate background rejections of misidentified hadronic decays provided by these working points in the interval $20 < p_T < 100$ GeV is 0.76, 0.17, 0.11 and 0.13% respectively. The reconstruction efficiency for medium muons which pass the trigger is above 99% for a wide p_T range as can be seen in Figure 4.11.

Muon isolation is defined using $p_T^{iso,var}$ as described in Section 4.3.3, but with the window extended to $\Delta R = \min(\frac{10\text{GeV}}{p_T[\text{GeV}]}, 0.3)$. A “FixedCutTightTrackOnly” isolation definition is used in this work, which applies a requirement $p_T^{iso,var}/p_T(\mu) < 0.06$, as well as “LooseTrackOnly” definition which is 99% efficient, constant in p_T and η [177].

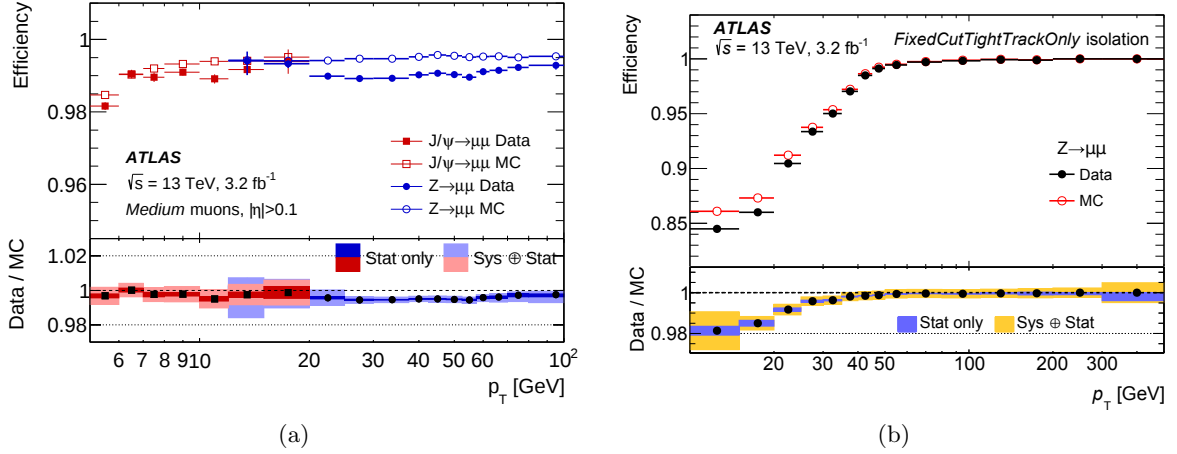


Figure 4.11: Efficiency of a) the medium identification working point and b) the FixedCut-TightTrackOnly isolation working point as a function of muon p_T [177]. Data (solid point) and MC (hollow point) are shown separately. The bottom subplot shows the ratio of data to MC, with the statistical uncertainty and total uncertainty shown as filled bands.

4.3.5 Missing Transverse Energy

Neutrinos are not expected to interact with the detector, but still may carry a significant momentum with respect to the rest of the event. Due to the effects of PDFs (see Section 4.1) the full energy and momentum conservation of an event can not be constrained as the exact parton momentum is unknown. The energy and momentum along just the transverse directions are still conserved and is effectively zero, as the incoming states are designed to be traveling nearly parallel to the beam-pipe near the collision. By measuring the energy imbalance in the transverse plane, some of the kinematic quantities related to neutrinos in the event can then be reconstructed. This energy imbalance is conventionally called missing transverse energy and labeled E_T^{miss} .

One subtlety that arises when defining E_T^{miss} is that it is specific to the exact definitions and calibrations of all physics objects in the event. The above reconstructed object definitions in Sections 4.3.1 - 4.3.4 are not orthogonal. In addition, definitions of τ -leptons [179, 180] and photons [181, 182], which are also needed for full event reconstruction also overlap with these definitions. To solve this overlap, a prioritization scheme is applied where objects lower on the list are removed from the E_T^{miss} calculation if they have overlapping calorimeter deposits or tracks with an object higher in priority. The ATLAS priority list is: electrons, photons, τ -leptons, muons, jets, and tracks associated to the primary vertex. The specific object definition choice and overlap procedure is detailed in Ref. [183].

The E_T^{miss} value for an event is then defined as the magnitude of the 2-vector given by [183,

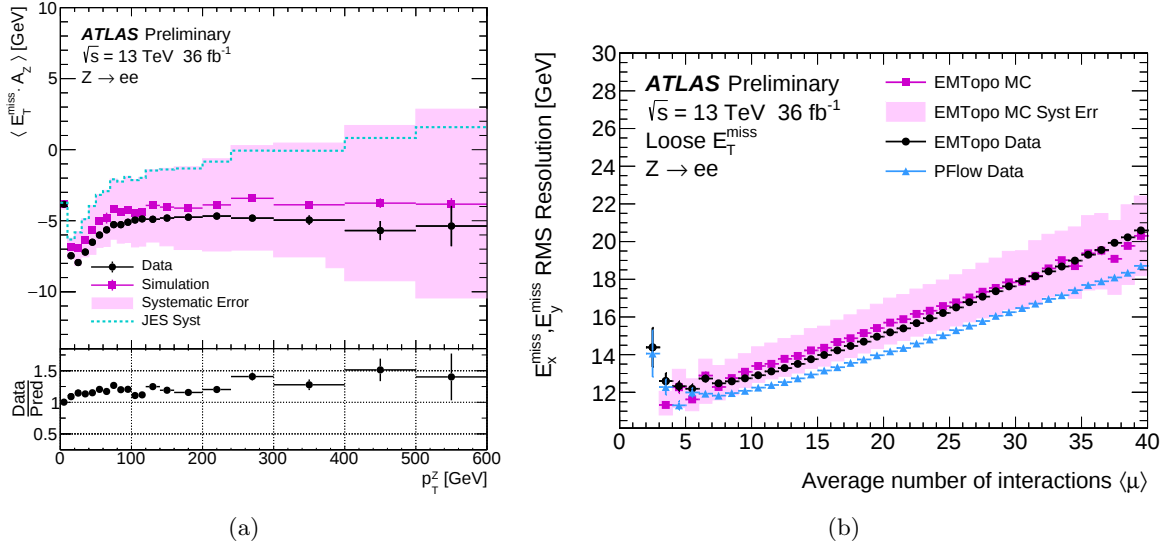


Figure 4.12: The a) effective E_T^{miss} scale offset with respect to prediction as a function of a reference object p_T , and b) the RMS of the E_T^{miss} variable as a function of $\langle \mu \rangle$. Both are estimates in $Z \rightarrow e\bar{e}$ events and the reference object is the reconstructed Z momentum [183].

184]

$$\vec{E}_T = \sum_e \vec{p}_T(e) + \sum_{\text{photons}} \vec{p}_T(\gamma) + \sum_{\text{tau}} \vec{p}_T(\tau) + \sum_{\text{muons}} \vec{p}_T(\mu) + \sum_{\text{jets}} \vec{p}_T(j) + \sum_{\text{tracks}} \vec{p}_T(trk) \quad (4.12)$$

where the sums run over all candidates which pass the object definition and overlap removal. The first five terms sum the mainly hard-particle contributions while the last track term is used to include the remaining soft radiation coming from the primary vertex.

The E_T^{miss} calibration is reliant on the calibration of all the individual objects in the calculation. The exact E_T^{miss} scale and resolution are topology dependent, especially on the contribution of jets and soft radiation in the event. The systematics associated to the E_T^{miss} scale and resolution from the hard-scatter contributions are propagated from the uncertainties from their dedicated calibrations. The uncertainty associated to the soft-track terms are measured $Z \rightarrow \ell\bar{\ell}$ in events [183, 184].

The scale offset and resolution of the E_T^{miss} variable as observed in data and MC for $Z \rightarrow e\bar{e}$ events is shown in Figure 4.12. As described above, distributions of these variables are topology dependent but agreement between data and MC is found within systematic uncertainties. The $Z \rightarrow ee$ topology mainly estimates the contributions from jets and soft-tracks, since the only other expected contribution is from final-state electrons, which are used as the reference probe for this study [183].

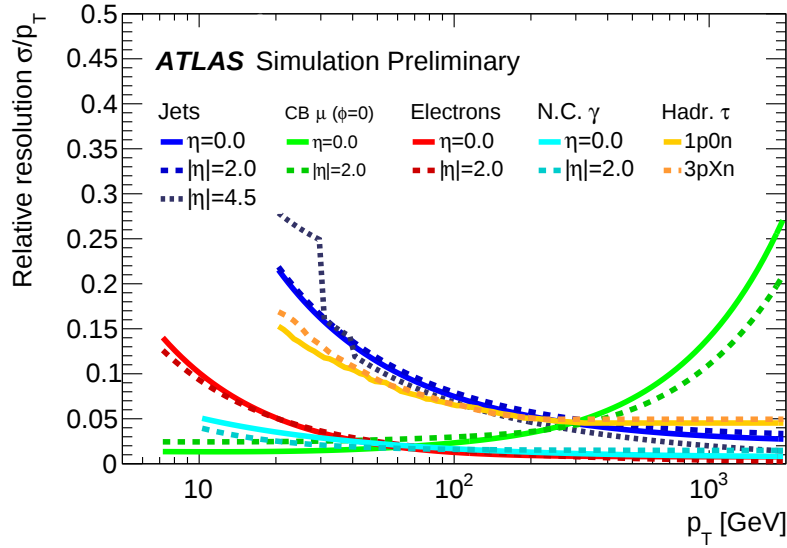


Figure 4.13: Relative p_T uncertainties for reconstructed small- R jets (blue), muons (green) and electrons (red) [185]. Similar curves for photon and hadronically decaying τ -leptons are shown but not discussed in this work.

Summary

Connecting detector observables to theoretical prediction is a two-fold process. First an understanding of the theoretical modeling and detector simulation must be known well in order to connect theory to experiment. Secondly, a clear algorithm is needed to relate detector level observables back to particles of the underlying theory.

This second step forms physics objects, which typically need to be calibrated to match truth-level predictions. Similarly, MC predictions need to be adjusted to properly represent data in areas where the MC modeling is not accurate. These calibration steps introduce several additional uncertainties on the reconstructed object quantities. All these steps are specific to the particular type of underlying particle which is attempted to be reconstructed. A comparison of the relative p_T uncertainties for several reconstructed object definitions can be seen in Figure 4.13.

Chapter 5

$VV \rightarrow \ell\nu qq$ Resonance Search with 36 fb^{-1}

Reports that say that something hasn't happened are always interesting to me, because as we know, there are known knowns; there are things we know we know. We also know there are known unknowns; that is to say we know there are some things we do not know. But there are also unknown unknowns—the ones we don't know we don't know.

—Donald Rumsfeld

As discussed in Section 2.3 many BSM models predict new particles which decay to pairs of electroweak vector bosons. The experimental signature of such processes would be a resonant structure in the invariant mass distribution of the diboson pair $m(VV)$, $V = W/Z$. Observation of such a resonant state would provide smoking-gun evidence for BSM physics as there is no SM process which could resonantly produce such a signature. The SM predicts a smoothly falling $m(VV)$ distribution from non-resonant diboson production and other SM processes which can be misidentified during reconstruction as diboson events.

W/Z bosons are not stable particles and decay into pairs of leptons or quarks. This poses a limitation as different final states typically require dedicated searches due to different background composition and experimental techniques required. For example in a pp collider, one of the dominant processes is inelastic proton scattering, which leads to a large number of multi-jet events ($O(10^6)$ pb for inclusive jet production), while the cross-section for SM events with leptons is orders of magnitude lower ($O(10^4)$ pb for inclusive W production) [186]. The expected background rates for final states with hadronically and leptonically decaying W/Z differ by orders of magnitude. Similarly, it is typically impractical to use MC to simulate multi-jet events, and hence data-driven techniques are usually used. For processes with leptons, which usually involve electroweak interactions, MC can be used to provide sufficiently accurate modeling.

Table 5.1 shows the branching fraction of different diboson states to purely quark final states (fully-hadronic), purely leptonic final states (fully-leptonic), and to states with one hadronically and one leptonically decaying boson (semi-leptonic). While the states with leptons are expected

to have smaller backgrounds, it can be seen that they also will have the smallest signal yields. For maximum single channel sensitivity, we will consider a search in semi-leptonic final states, which will provide a compromise between low backgrounds and high signal yields. We will find explicitly in Chapter 6 that the semi-leptonic channels are indeed one of the most sensitive channels for extracting limits on diboson resonances, notably in the TeV range.

The treatment of τ -leptons induces further complication as they decay promptly to states with a lighter charged lepton or pions. The reconstruction of the hadronic τ -lepton decays requires dedicated techniques which are not detailed in this work, and the leptonic decays can be difficult to distinguish from typical electron and muon production. We will not explicitly attempt to reconstruct τ -decays in this work, but will include them in background and signal calculations to account for the 34% branching fraction to lighter leptons. States with τ -leptons have been separated from the “leptonic decays” of Table 5.1 and will not be considered explicitly when discussing “charged leptons” for the rest of the chapter.

This chapter will focus on the search for diboson channels in the final state with one leptonic W and one hadronic W/Z , known as the $\ell\nu qq$ semi-leptonic final state¹, with 36 fb^{-1} of ATLAS data as published in Ref. [8]. Preceding this publication, similar searches had been done with lower center of mass energies and integrated luminosities by both the ATLAS [187–189] and CMS [190] experiments. In comparison to previous publications, this publication includes a reoptimized channel for low-mass resonances as well as a new interpretation in the context of vector-boson fusion (VBF) produced signals. Diboson resonance searches in other final states have also been published by the ATLAS collaboration with the same dataset [191–195].

The analysis is mainly designed around predictions from MC, which are described in Section 5.1 for the background and signals. An intricate event selection criteria designed to optimize the search sensitivity is discussed in Section 5.2. Multiple signal regions are developed to select events where the quarks from the hadronically decaying boson can be individually reconstructed or overlap. Further regions are designed to enhance the sensitivity to signals produced via VBF by requirements on additional energetic and forward hadronic radiation in the event, which is not expected to be present in $ggF/q\bar{q}$ produced signals. Control regions, very similar to the main signal regions but with inverted cuts are also defined to help constrain the normalization on the main background processes. These also constrain the large number of systematic uncertainties described in Section 5.3.

The final observable used in the analysis is the reconstructed diboson invariant mass $m(WV)$. The binning is chosen such that the narrow resonances probed in this work should appear as an excess of data events with respect to the SM prediction in one or two bins. The $m(WV)$ distribution of all the signal and control regions are fitted simultaneously using a profile-likelihood approach detailed in Appendix A and the final results shown in Section 5.4.

¹For labeling of final states, the anti-particle label will be removed. We will also denote $W \rightarrow \ell\nu$ to represent $W^- \rightarrow \ell\bar{\nu}$ and $W^+ \rightarrow \bar{\ell}\nu$. This convention is also used in Chapters 6 - 8.

$\ell = e, \mu, \nu$	WW	WZ	ZZ
$\ell\ell\ell\ell$	4.7%	5.8%	7.1%
$\ell\bar{\ell}q\bar{q}$	29.2%	33.1%	37.4%
$q\bar{q}q\bar{q}$	45.7%	46.8%	48.8%
taus	20.4%	14.3%	6.7%

Table 5.1: The branching fraction of pairs of different diboson states to combinations of hadronic decays $q\bar{q}$ and leptonic decays $\ell\bar{\ell}$. The leptonic decays do not include states with τ -leptons which have been grouped into a separate “taus” category.

5.1 Signal and Background Simulation

This analysis will rely mainly on predictions from MC generators to model both the expected signal and background. The MC predictions will be used to optimize the selection criteria as well as provide templates for the expected background distributions in the final fit to data. The choice of generator for each sample is taken to be the one found to best represent ATLAS data in dedicated studies [196–199].

5.1.1 Signal Models

Several models predict new resonances which decay to the diboson $\ell\nu qq$ state as detailed in Section 2.3. In the models discussed there, the theoretically motivated resonance masses span the range of several hundred GeV up to several TeV. To cover a wide range of models but still remain as inclusive as possible, the results of this analysis will be interpreted in three specific benchmark models assuming different spin hypothesis for the new resonant particle: spin-0, 1, or 2. For all benchmark models under consideration the intrinsic width of the resonance is narrower than the detector resolution.

Neutral and charged spin-1 resonances are interpreted in the context of a Heavy Vector Triplet (HVT) Z' and W' described in Section 2.3.2. Spin-2 resonances, which are predicted in models of quantum gravity, are interpreted in the Bulk Randall-Sundrum (RS) model of Section 2.3.3 as a Kaluza-Klein graviton G_{KK} . Lastly, spin-0 resonances are probed in a naive heavy Higgs-like scalar model where the new resonance H has Higgs-like coupling and negligible width. Separate interpretations will be conducted for the heavy Higgs-like scalar and HVT model under gluon-gluon fusion/quark-anti-quark annihilation ($ggF/q\bar{q}$) production and vector-boson fusion (VBF) production, while the Bulk RS graviton will only be interpreted in the ggF production mode.

The Bulk RS graviton and HVT W'/Z' samples are both generated with MC@NLO 2.2.2 [130] with the NNPDF23LO [122] PDF set interfaced to Pythia 6 [200]. The Bulk RS graviton $G_{KK} \rightarrow WW$ samples are generated in the mass range of 300 - 5000 GeV with the model parameter $k/M_{Pl} = 1$ as detailed in Section 2.3.3. The HVT signals are generated in the same mass range for both $W' \rightarrow WZ$ and $Z' \rightarrow WW$ via $q\bar{q}$ and VBF production. The ggF samples

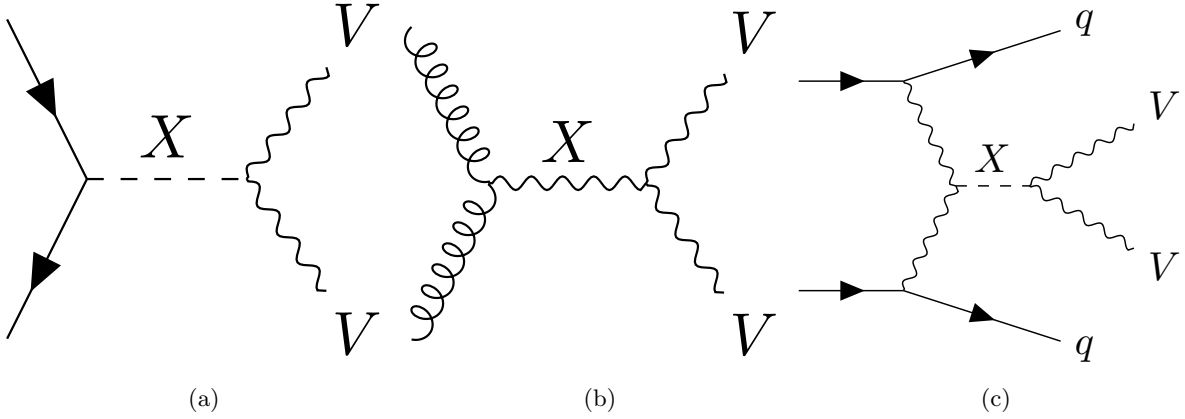


Figure 5.1: Representative Feynman diagrams for new resonances produced via a) $q\bar{q}$ annihilation, b) ggF , and c) VBF.

are produced with couplings set to the Model A benchmark, while the VBF ones use the Model C configuration where the fermion couplings are set to zero as detailed in Section 2.3.2. The Bulk RS graviton is produced only in the ggF production mode. Feynman diagrams illustrating the different production modes for a possible resonant state X are shown in Figure 5.1.

The heavy Higgs-like scalar samples decaying as $H \rightarrow WW$, are generated with POWHEG-Box v1 [201] with dedicated configuration for NLO calculations of the Higgs in ggF and VBF production [202, 203]. The showering and hadronization is done with Pythia 8 [133]. To simulate the narrow width approximation, the scalar width is set to 4 MeV for all mass points generated in the interval of 300 - 3000 GeV.

5.1.2 Background Processes

There are four SM processes which are expected to form the majority of the background for this search in the $\ell\nu qq$ final state: V +jets, $t\bar{t}$, single-top, and SM VV . We will consider W +jets and Z +jets separately, as the latter are sub-dominant background contribution for this analysis. The majority of the backgrounds are reducible in the sense that they produce event topologies that mimic the diboson signal, but do not actually contain two W/Z bosons, and hence cuts can be applied to reduce their contributions.

The main background originates from W +jets, where there is a W boson decaying leptonically which is recoiling off additional jet radiation. The leptonic decay products of the W can pass the trigger and identification requirements while the additional jets can be misidentified as the hadronically decaying boson. In comparison to the signal process, the jets in the V +jets events will have a diffuse spread in the invariant mass of the jets as opposed to a peaked distribution at the W/Z mass. Z +jets is expected to contribute significantly less than W +jets to the background, as a strict requirement on the charged lepton multiplicity will suppress both $Z \rightarrow \ell\bar{\ell}$ and $Z \rightarrow \nu\bar{\nu}$ contributions. As discussed in Sections 4.3.3 - 4.3.4, the charged lepton

efficiencies are greater than 80% and the misidentification rates below 1%, which suppresses the Z +jets contributions.

The next major source of background is $t\bar{t}$ events with one semileptonic top decay. In such processes there are two W bosons, one decaying leptonically and the other hadronically, which naively look like the signal. These events can be differentiated with respect to the signal due to the additional jets from the top decay, which are expected to be b -tagged. The fully leptonic and hadronic $t\bar{t}$ events are not expected to contribute significantly as the majority of these events can be removed by cuts on the lepton multiplicity. The production of single-top events can also contribute to our background, 90% of which is expected to be from the Wt -channel. Similar to $t\bar{t}$, the single-top events are expected to have b -tagged jets which can be used to separate them with respect to the signal.

The last sub-dominant contribution to be considered is the non-resonant diboson (SM VV) production. This is an irreducible background, which only differs from the signal in the sense that it is non-resonant.

To provide intuition on the scale of these backgrounds, the inclusive cross-sections of these processes as measured by the ATLAS experiments are approximately: $2 \times 10^5 \text{ pb}$ for V +jets, 800 pb for $t\bar{t}$, 300 pb for single-top, and 200 pb for SM VV [186]. As can be seen, W +jets provides the highest expected background yields followed by $t\bar{t}$. This is only for illustrative purposes though, as each background has different kinematic distributions and hence have different efficiencies to pass the event selection described in Section 5.2.

The V +jets samples are generated using Sherpa v2.2.1 [204] with the NNPDF30NNLO set [121]. Matrix elements were calculated for up to 2 partons at NLO and 4 partons at LO using Comix [205] and OpenLoops [206] and merged with the Sherpa parton shower [207] using the ME+PS@NLO prescription [208]. To provide sufficient statistics over a large region of phase space, in particular in the high energy tails, the samples are sliced by $\max(H_T, p_T(V))$ ranges, where H_T is the scalar p_T sum of all jets in the event. These samples are further filtered into samples with b -quarks, c -quarks, or solely light flavor jets in the final state. The overall cross-section includes up to NNLO QCD corrections [209]. SM diboson production are modeled with a similar setup as the V +jets samples, but calculated to NLO for only one additional parton and LO for up to 4.

Backgrounds with top quarks are simulated with the POWHEG-Box v2 [201] matrix-element calculation with CT10 PDF [125] interfaced to Pythia 6.426 [200]. These include $t\bar{t}$ and single-top in the s and Wt -channels. The t -channel single-top production is modeled with POWHEG-Box v1 [201]. For all top processes the top-quark spin correlations are preserved and the top quark mass was set to 172.5 GeV. The EvtGen v1.2.0 [135] code is used to improve the simulation of heavy flavour hadron decays. The $t\bar{t}$ cross section is calculated at NNLO in QCD, including resummation of next-to-next-to-leading logarithmic (NNLL) soft-gluon terms [210, 211]. The single-top cross sections are calculated to NLO QCD [212], with additional corrections in the Wt -channel for the soft-gluon resummation at NNLL [213].

For all samples the effect of pile-up is simulated by overlaying additional inelastic pp collisions, with the distribution of the number of overlaid events set to match pile-up conditions in data. These additional events are generated with Pythia v8.186 [200] and the MSTW2008LO [214] PDF set. For all signal models EvtGen v1.2.0 [135] is used for b and c -hadron decays. All samples include τ leptons in W/Z decays.

5.2 Event Selection

The main analysis strategy involves the search for resonances in the reconstructed invariant mass distribution of events consistent with a leptonically decaying W and a hadronically decaying W/Z . Events are required to have either a single charged lepton (an electron or muon), E_T^{miss} from a neutrino, and jets from the hadronic $V = W/Z$ decay. After the object-based event selection, an additional set of kinematic cuts are applied to reduce the background contribution from events with misidentified V bosons.

One complication which arises is the reconstruction of the $V \rightarrow q\bar{q}$ decay over a wide p_T range. As discussed in Section 4.3.2 the opening angle between two outgoing particles in the 2-body decay of a boson follows Equation 4.5 which is restated here

$$\Delta R \approx \frac{2m(V)}{p_T(V)}$$

where $m(V)$ and $p_T(V)$ are the mass and p_T of the parent boson. For W/Z bosons with $p_T \lesssim 450\text{GeV}$, the opening angle is sufficiently large that the two outgoing quarks can be reconstructed as individual small- R jets. Above this range the radiation patterns overlap and reconstitution of the two individual decays becomes difficult. Instead, it becomes more efficient to attempt to reconstruct such decays as a single large- R jet. To provide optimal sensitivity in the whole p_T range, and hence mass range of the new resonant particle, the search will be divided into two selections optimized for decays which have two small- R jets, known as the resolved channel, or a single large- R jet, known as the merged channel.

The selection will also be divided to target signals produced in the $ggF/q\bar{q}$ or VBF topology. Signals produced via the VBF process are predicted to have additional high- p_T forward jets which are not present in the $ggF/q\bar{q}$ topologies. The VBF channel will have additional requirements on the jets in the event to provide a low-background channel with increased sensitivity to VBF signals. Events which do not pass the VBF channel requirements will be identified as candidate $ggF/q\bar{q}$ events.

The selection common to all analysis regions is discussed in Section 5.2.1 and the categorization of VBF and $ggF/q\bar{q}$ events is detailed in Section 5.2.2. The selection of events in the merged channel is detailed in Section 5.2.3 and the resolved channel selection in Section 5.2.4. The merged selection will be subdivided into a high-purity (HP) and low-purity (LP) channels to optimize the signal acceptance and sensitivity.

In total there are six signal regions (SR), for merged HP, merged LP, and resolved searches in each of the ggF/ $q\bar{q}$ and VBF regions. Additional control regions will be identified as sideband regions in Section 5.2.5 which mimic the respective signal regions, but fail a certain cut definition. These regions will be used to constrain the normalization of the major backgrounds in a data-driven way. Explicitly, W +jets control regions (WCR) will be designated for events which fail requirements on the $V \rightarrow q\bar{q}$ invariant mass, and $t\bar{t}$ control regions (TCR) will gather events which fail the requirements on the number of b -tagged jets in the event. Lastly Section 5.2.6 will detail an orthogonality procedure to simplify the statistical analysis by requiring events to be categorized into only a single region. Summary tables of the selection criteria in the merged and resolved regions can be found in Tables 5.2 - 5.3.

The final fits to data, described in Section 5.4, will be separated into ggF/ $q\bar{q}$ and VBF channels as well as WW and WZ channels. For each of the ggF/ $q\bar{q}$ and VBF fits, the three relevant signal and six control regions are fit simultaneously. The WW and WZ interpretations only differ by whether the events in the signal region explicitly pass the W or Z mass requirements on the selected jets.

5.2.1 Common Selection

The common event and trigger selection is oriented around the $W \rightarrow \ell\nu$ selection, which does not differ much between the various signal regions. Each event is required to have a single “tight” electron or “medium” muon with $p_T > 27 \text{ GeV}$, denoted the signal lepton, and no additional “loose” electrons or muons with $p_T > 7 \text{ GeV}$ (see Sections 4.3.3 - 4.3.4 for electron and muon definitions). This suppresses the majority of the multi-jet backgrounds and electroweak processes which have additional charged leptons. The identification definitions include the criteria that medium (loose) muons are within $|\eta| < 2.5$ (2.7) and electrons are within $|\eta| < 2.47$. Electrons are required to not be within the calorimeter crack region of $1.37 < |\eta| < 1.52$ where energy deposits are not measured as accurately. These identification requirements are expected to be 88% and 99% efficient for real electrons and muons at $p_T > 100 \text{ GeV}$. Furthermore, the ID tracks associated to the charged leptons are required to satisfy $|d_0/\sigma(d_0)| < 5$ (3) and $|z_0 \times \sin \theta| < 0.5 \text{ mm}$ for electrons (muons). The “FixedCutTight” and “FixedCutTightTrackOnly” isolation requirements as detailed in Sections 4.3.3 - 4.3.4 are applied to signal leptons and the “LooseTrackOnly” requirement is places on the loose charged leptons.

The p_T and ϕ of the neutrino is reconstructed directly from the $\vec{E}_T$ calculated in the event (see Section 4.3.5). The full 4-momenta of the neutrino in the event is reconstructed by assuming the invariant mass of the lepton-neutrino pair is consistent with the W mass. This provides a quadratic equation for the third component of the neutrino momenta p_z , and the solution with the smallest real part is chosen. This procedure allows for better resolution as the diboson invariant mass $m(WZ)$ can be fully reconstructed. The invariant mass distribution has a full-width half-max a factor of 2 smaller than the transverse mass defined as $m_T(WZ) =$

$m^2(J\ell) + 2 \left(E_T(J\ell) \times E_T^{\text{miss}} - \vec{p}_T(J\ell) \cdot \vec{E}_T^{\text{miss}} \right)$, which makes no use of the neutrino p_z component.

Events are recorded from data using a fast hardware and software trigger system detailed in Section 3.2.5. For consistency, events considered in this analysis are required to pass trigger requirements which match the lepton content expected for the signals. Events are triggered using the lowest threshold unprescaled lepton and E_T^{miss} triggers available during that operations period, which are explicitly detailed in Section 3.2.5. For single lepton triggers, the selected offline lepton is required to match the online triggered lepton. The offline lepton p_T thresholds are chosen so that selected events are above the trigger plateau threshold.

5.2.2 VBF Channels

The cross-sections for VBF produced signals can be orders of magnitude less than those produced by $ggF/q\bar{q}$ production, dependent on the signal model. To improve the sensitivity for the VBF signals, a dedicated VBF channel is developed which exploits the unique topology of VBF events to reduce backgrounds. Events with a resonant state produced via the VBF production mechanism are expected to have two additional forward jets which are not present in the $ggF/q\bar{q}$ topology or many backgrounds. The first stage of the analysis after the trigger requirement and the selection of leptons is to categorize events into separate $ggF/q\bar{q}$ and VBF channels based on whether they have jets consistent with the VBF topology. If an event passes this requirement, it is considered a candidate VBF event, and if not it is candidate $ggF/q\bar{q}$ event. The remainder of the selection for each category is almost identical, with separate merged and resolved channels for both.

Only events which have two small- R jets in addition to a large- R jet, or events with greater than four small- R jets are considered as possible VBF events (jet collections are defined in Section 4.3.2). For these events, the two small- R jets in the event with the highest invariant mass $m(j_1^{VBF}, j_2^{VBF})$ which satisfy, $p_T > 30$ GeV, $|\eta| < 4.5$, $\eta(j_1^{VBF}) \cdot \eta(j_2^{VBF}) < 0$ and are not b -tagged are selected as the VBF-tag jets. These requirements are applied since VBF-tag jets are usually highly energetic, in opposite hemispheres and relatively forward in the detectors.

To further reduce backgrounds, additional requirements are placed on the VBF-tag jets in order to be classified as a candidate VBF event. For signal events, the VBF-tag jets are expected to become increasingly energetic and spatially separated as the resonant mass increases. Two cuts on the VBF-tag jet invariant mass $m(j_1^{VBF}, j_2^{VBF})$ and the relative opening angle $\Delta\eta(j_1^{VBF}, j_2^{VBF})$ are optimized for maximum signal sensitivity to low-mass signals under 1 TeV. Higher mass signals are expected to pass all cuts derived this way. The metric used for the optimization is the binned significance calculated as

$$\sigma = \sqrt{\sum_i^N \left(\frac{s_i}{\sqrt{s_i + b_i + (\Delta \cdot b_i)^2}} \right)^2} \quad (5.1)$$

where s_i and b_i are the expected signal and background yields in each bin i of the $m(WV)$

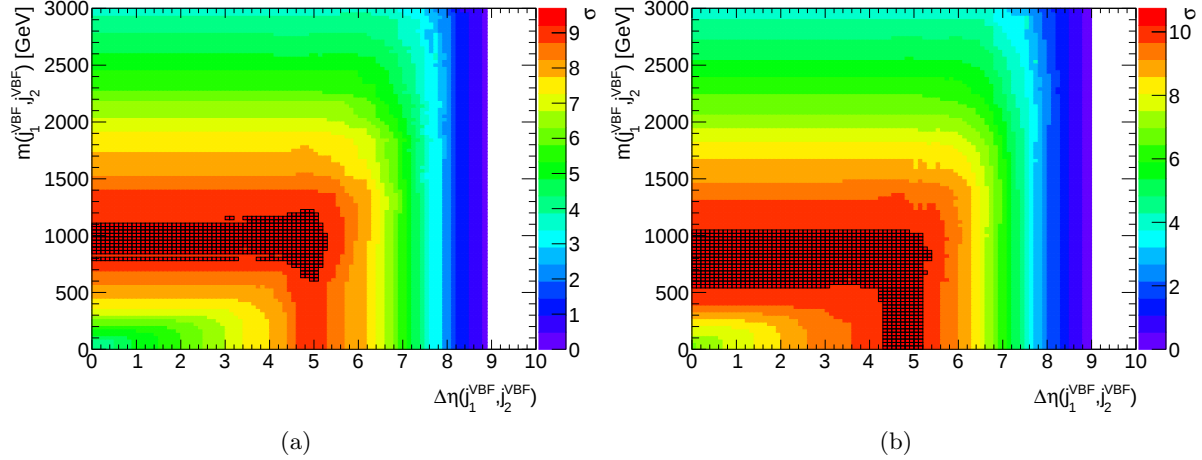


Figure 5.2: The signal significance as a function of cuts on the two VBF-tag jet invariant mass $m(j_1^{VBF}, j_2^{VBF})$ and η opening $\Delta\eta(j_1^{VBF}, j_2^{VBF})$ for a 500 GeV and 700 GeV heavy Higgs-like scalar signal. The region highlighted by black squares shows the cut range where the signal significance is within 5% of the maximum. The magnitude of the significance is model dependent, but the relative differences along the cuts is not.

distribution. A background uncertainty rate of $\Delta = 0.015$ is used in the calculation to include systematic effects in the calculation. Figure 5.2 shows the binned significance as a function of the cut value for 500 GeV and 700 GeV signal masses. The area highlighted by black boxes is the region which varies from the maximum by less than 5%. The distributions show that the significance is flat around the peak values, hence the signal significance is not sensitive to small changes in the cut values. The final cut values are chosen to be $\Delta\eta(j_1^{VBF}, j_2^{VBF}) > 4.7$ and $m(j_1^{VBF}, j_2^{VBF}) > 770$ GeV, for consistency with the $\ell\ell qq + \nu\nu qq$ publication [191]. This selection is approximately 28% efficient for VBF signals, but reduces the background composition by two orders of magnitude. Candidate VBF events not passing the above selection are recycled into the ggF/ $q\bar{q}$ selection.

Events which pass the VBF selection, have the two VBF-tag jets removed from the possible collection of $V \rightarrow q\bar{q}$ candidates in the resolved channels, and in the merged channel the large- R jets are required to have $\Delta R(J, j_i^{VBF}) > 1.5$. If the events fall into the ggF/ $q\bar{q}$ category, the VBF-tag jets designation is removed, and all jets are considered for the respective merged and resolved selections.

5.2.3 Merged Selection

The merged selection, where the $V \rightarrow q\bar{q}$ is reconstructed as one large- R jet, is optimized to identify signals with resonance mass $m > 1$ TeV. The full kinematic cut requirements for this channel are shown in Table 5.2 and are detailed further below.

To pass this selection, there must exist one large- R jet within $|\eta| < 2.5$ which has $p_T > 200$ GeV and $m > 50$ GeV. These jets are groomed using a trimming procedure and all mass

Selection		SR: HP (LP)	W CR: HP (LP)	$t\bar{t}$ CR: HP (LP)
Production category	VBF	$m^{\text{tag}}(j, j) > 770 \text{ GeV}$ and $ \Delta\eta^{\text{tag}}(j, j) > 4.7$		
	ggF/ $q\bar{q}$	Fails VBF selection		
$W \rightarrow \ell\nu$ selection	Num. of signal leptons	1		
	Num. of veto leptons	0		
	$E_{\text{T}}^{\text{miss}}$	$> 100 \text{ GeV}$		
	$p_{\text{T}}(\ell\nu)$	$> 200 \text{ GeV}$		
	$E_{\text{T}}^{\text{miss}}/p_{\text{T}}(e\nu)$	> 0.2		
$V \rightarrow J$ selection	Num. of large- R jets	≥ 1		
	D_2 eff. working point (%)	Pass 50 (80)	Pass 50 (80)	Pass 50 (80)
	Mass window Eff. working point (%)	Pass 50 (80)	Fail 80 (80)	Pass 50 (80)
Topology criteria	$p_{\text{T}}(\ell\nu)/m(WV)$ $p_{\text{T}}(J)/m(WV)$	> 0.3 for VBF and > 0.4 for ggF/ $q\bar{q}$ category		
Num. of b -tagged jet	excluding b -tagged jets with $\Delta R(J, b) \leq 1.0$	0		≥ 1

Table 5.2: Selection requirements for the merged $\ell\nu q\bar{q}$ signal and control regions. Further details in text.

values are calculated with the combined mass definition (detailed in Section 4.3.2).

The requirements on the reconstructed leptonic $W \rightarrow \ell\nu$ are tightened to match the fact the expected W and hadronic V are expected to both carry approximately equal momenta if produced by an on-shell resonance. These tighter requirements are $E_{\text{T}}^{\text{miss}} > 100 \text{ GeV}$, and $p_{\text{T}}(\ell\nu) > 200 \text{ GeV}$. For the same reason, cuts on $p_{\text{T}}(W/V)/m(WV)$ are expected to increase the background rejection. This cut is separately optimized for the ggF/ $q\bar{q}$ and VBF topologies as some model dependency is found in the signal distributions. The threshold is chosen to be $p_{\text{T}}(W/V)/m(WV) > 0.4$ for the ggF/ $q\bar{q}$ channel and 0.3 for the VBF channel. An additional cut imposed only in the electron channel is that $E_{\text{T}}^{\text{miss}}/p_{\text{T}}(e\nu) > 0.2$. Such a cut removes a significant fraction of the remaining multi-jet background and only reduces the signal acceptance at 3 TeV by approximately 3%. After this cut the multi-jet background is expected to be negligible in the electron channel. No additional cuts in the muon channel are needed to suppress multi-jet events to a negligible value.

To reduce the contribution from $t\bar{t}$ events, events are required to have no additional b -tagged jets $\Delta R(J, b) > 1$ from the large- R jet. This criterion allows for b -jets within the large- R jet in order to not suppress the $Z \rightarrow b\bar{b}$ decays. After applying this cut, the number of $t\bar{t}$ events is reduced by approximately 70%.

The majority of the discrimination power in the merged channel comes from exploiting the substructure found in boosted $V \rightarrow q\bar{q}$ decays. It has been found in several ATLAS publications that the large- R jet mass and the $D_2^{\beta=1}$ variable are strong discriminants between hadronically decaying bosons and multi-jets events [161, 164, 167]. In this analysis, the distribution of these variables in the selected V +jets background events are expected to be similar to the multi-jets in Figure 4.8. This is due to the stringent requirements on the leptonic boson decay, which then require the large- R jet in V +jets events to come from additional QCD radiation. To provide

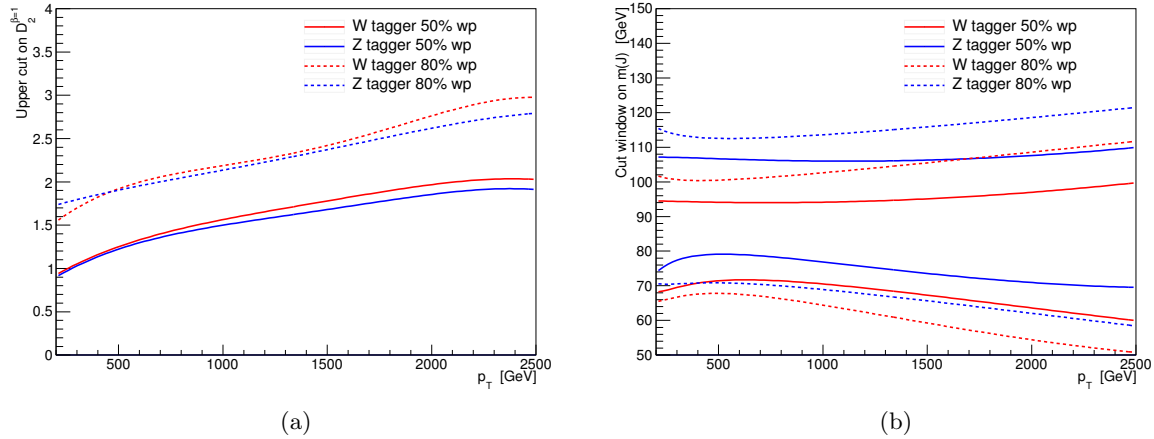


Figure 5.3: a) Upper cut on $D_2^{\beta=1}$ and b) mass window used as a function of p_T for the W -tagger (red) and Z -tagger (blue) for fixed efficiency working points of 50% (solid) and 80% (dashed).

increased signal to background separation, a hadronic V -tagger was designed using these two variables. The V -tagger was optimized such that p_T -dependent cuts on $D_2^{\beta=1}$ and the jet mass are derived which provide uniform efficiency working points for HVT signals and optimal background rejection in each p_T bin. A W -tagger and Z -tagger were optimized independently for two working point with flat 50% and 80% efficiency. The cut boundaries used by the V -tagger as function of p_T can be found in Figure 5.3. The main difference between the W -tagger and Z -tagger is that the mass windows are centered approximately at the W mass and Z mass respectively. The background rejection of the taggers vary from 30 - 70 dependent on the jet p_T .

An important point to note, is the HVT signal produces longitudinal gauge bosons and so the V -tagger may be sub-optimal for transversely polarized bosons. All benchmark models considered in this analysis produce only longitudinal polarized bosons, so this subtlety is not relevant for these models. When re-interpreting the results of these benchmark signals to other models with transverse bosons, care should be taken to account for this effect.

At this point the merged channels are further divided into High-Purity (HP) and Low-Purity (LP) selections. Events with large- R jets that pass the 50% V -tagging working point are grouped into the HP channel. To recover signal acceptance, events which fail the 50% working point but pass the 80% working point are grouped into the LP channel. Recovering these events increases the overall signal significance by more than 30% for signal masses above 3 TeV.

After all selections, the background composition in the $ggF/q\bar{q}$ HP (LP) region is 55 (72)% W +jets and 32 (19)% $t\bar{t}$. These values in the VBF HP (LP) region are 40 (60)% W +jets and 50 (30)% $t\bar{t}$. The remaining contributions come from single-top, SM VV , and Z +jets events.

Selection		WW (WZ) SR		W CR	t $\bar{t}$ CR
Production category	VBF	$m^{\text{tag}}(j, j) > 770 \text{ GeV}$ and $ \Delta\eta^{\text{tag}}(j, j) > 4.7$			
	ggF/q $\bar{q}$	Fails VBF selection			
$W \rightarrow \ell\nu$ selection	Num. of signal leptons	1			
	Num. of veto leptons	0			
	$E_{\text{T}}^{\text{miss}}$	$> 60 \text{ GeV}$			
	$p_{\text{T}}(\ell\nu)$	$> 75 \text{ GeV}$			
	$E_{\text{T}}^{\text{miss}}/p_{\text{T}}(e\nu)$	> 0.2			
$V \rightarrow jj$ selection	Num. of small- R jets	≥ 2			
	$p_{\text{T}}(j_1)$	$> 60 \text{ GeV}$			
	$p_{\text{T}}(j_2)$	$> 45 \text{ GeV}$			
	$m(jj)$ [GeV]	[66, 94] ([82, 106])	< 66 or [106, 200]	[66, 106]	
Topology criteria	$\Delta\phi(j, \ell)$	> 1.0			
	$\Delta\phi(j, E_{\text{T}}^{\text{miss}})$	> 1.0			
	$\Delta\phi(j, j)$	< 1.5			
	$\Delta\phi(\ell, E_{\text{T}}^{\text{miss}})$	< 1.5			
	$p_{\text{T}}(\ell\nu)/m(WV)$ $p_{\text{T}}(jj)/m(WV)$	> 0.3 for VBF and 0.35 for ggF/q $\bar{q}$ category			
	Num. of b -tagged jets	$j_1 \equiv b$ or $j_2 \equiv b$ where $V \rightarrow j_1 j_2$	$\leq 1(2)$	≤ 1	> 0 (for jets other than j_1 or j_2)
$j_1 \neq b$ and $j_2 \neq b$ where $V \rightarrow j_1 j_2$		0			

Table 5.3: Selection requirements for the resolved $\ell\nu qq$ signal and control regions. Further details in text.

5.2.4 Resolved Channel

The resolved channel is optimized for signal acceptance for resonances with mass $< 1\text{TeV}$. The full selection requirements are shown in Figure 5.3. In the resolved selection, the lepton requirements are increased on the $E_{\text{T}}^{\text{miss}}$ and $p_{\text{T}}(\ell\nu)$ to 60 GeV and 75 GeV respectively. The requirement thresholds are lower than that corresponding values in the merged selection. Requirements on the $p_{\text{T}}(W/V)/m(WV)$ value in a resolved event are optimized in the same way as in the merged channel. The optimal values for the relative momentum cut in the VBF channel is found to be the same and a lower threshold of 0.35 is found in the ggF/ $q\bar{q}$ channel. The same $E_{\text{T}}^{\text{miss}}/p_{\text{T}}(e\nu)$ cut as the merged analysis is kept to suppress multi-jet contributions.

In the resolved selection, the $V \rightarrow q\bar{q}$ decay is reconstructed as the two small- R jets with highest p_{T} . If a candidate $V \rightarrow q\bar{q}$ jet pair is found, the leading jet is required to have $p_{\text{T}} > 60 \text{ GeV}$ and the sub-leading jet $p_{\text{T}} > 45 \text{ GeV}$ to suppress backgrounds. Upper and lower mass window cuts are optimized using the binned significance definition in Equation 5.1 and with same procedure as implemented in Section 5.2.2. This procedure results in optimized windows of $66 < m(j_1, j_2) < 94 \text{ GeV}$ for $W \rightarrow q\bar{q}$ signals and $82 < m(j_1, j_2) < 106 \text{ GeV}$ for $Z \rightarrow q\bar{q}$. To reduce $t\bar{t}$ contamination, two of the jets are allowed to be b -tagged in the Z selection. Including such events increases backgrounds by 5% but increases the signal acceptance by 15%

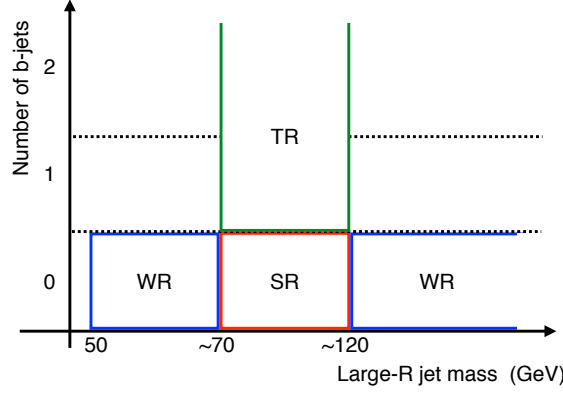


Figure 5.4: Sketch of the control region definition in the merged selection. The top control region has an inverted criteria on the number of b -tagged jets and the W +jets control region has inverted mass requirements (approximate window thresholds shown).

for Z events. Further suppression of $t\bar{t}$ events is accomplished by removing events with any additional jets which are b -tagged.

Additional topological cuts on the relative ϕ angles between the final four objects are applied. They are $\Delta\phi(j_i, \ell/\nu) > 1$, $\Delta\phi(j_1, j_2) < 1.5$ and $\Delta\phi(\ell, \nu) < 1.5$. These cuts are fully efficient for the signals considered. After all selection requirements, the expected background events in the $ggF/q\bar{q}$ (VBF) region are 67 (60)% W +jets and 24 (30)% $t\bar{t}$. To restate for completeness, the remaining background events are single-top, SM VV , and Z +jets.

For the final $m(WV)$ distribution, it is found that the signal resolution can be improved by 20% if the $V \rightarrow q\bar{q}$ dijet system is constrained to the W or Z mass (m_W or m_Z). After the selection described above, the dijet system 4-momenta is adjusted to

$$p_T^{corr}(jj) = p_T(jj) \frac{m(jj)}{m_{W/Z}} \quad (5.2)$$

$$m^{corr}(jj) = m_{W/Z}$$

where $m_{W/Z}$ is either m_W or m_Z when interpreting the results for signals with WW or WZ decays. Application of this procedure on the background samples shows no noticeable sculpting of the background $m(WZ)$ distribution.

5.2.5 Background Determination

In each of the signal regions, greater than 80% of the total background is from W +jets and $t\bar{t}$ events. An accurate measurement of these contributions is then important for a reliable estimation of the background. This is especially true for the overall normalizations of both processes as it is not expected that the inclusive MC cross-sections used to generate the samples will be appropriate after the exclusive requirements of the signal regions.

Dedicated background estimations are done with control regions defined to match as kinematically close as possible the signal regions (SR), but contain high purities of W +jets and $t\bar{t}$ events. For each of these control regions the selection will be identical except for one requirement which will be inverted. The definition of the control regions for the merged channels is sketched in Figure 5.4. Top control regions (TCR) are defined to have the same selection as the signal regions but with at least one additional b -tagged jet in the event. Separate TCR are defined for the merged HP, merged LP, and resolved regions. Similarly, the W +jets control regions (WCR) match the signal region definitions except the requirements on mass windows are inverted. For the merged region this requires that the large- R fails the 80% working point mass windows for both W and Z taggers. The merged HP WCR and LP WCR are distinguished if the large- R jet passes the 50% $D_2^{\beta=1}$ cut or fails the 50% but passes the 80% cut. In the resolved channel the WCR is defined such that the $V \rightarrow q\bar{q}$ jets must have either $m(j_1, j_2) < 66$ GeV or $106 < m(j_1, j_2) < 200$ GeV. The upper limit on the high-mass sideband is imposed such that equal events are expected in the upper and lower sidebands, in order to not bias the measurement.

The purity of the WCR is 65, 68, and 77% for the merged HP, merged LP and resolved selection. The same values are 85, 80, and 86% for the TCR. Even though the control regions are not 100% pure, the remaining contributions are mostly cross-contamination of $t\bar{t}$ in the WCR and W +jets in the TCR. An accurate measurement of the normalizations can then still be extracted with a simultaneous fit to both regions.

The remaining backgrounds are taken directly from MC predictions except for the multi-jet contribution. In the merged channel this contribution is expected to be negligible, while in the resolved region it is expected to contribute on the order of a few percent to the total background yield. The multi-jet contribution is estimated in the resolved regions using a data-driven “fake-factor method” as was done for the SM $VH \rightarrow \ell\nu b\bar{b}$ measurement by the ATLAS experiment [215].

5.2.6 Event Orthogonality Strategy

The majority of the signal and control region definitions above are designed to be orthogonal in the sense that no event can contribute to multiple regions. This is helpful as these regions can be statistically combined without the additional difficulty of measuring and applying correlations between overlapping events.

This is by design true between the ggF/ $q\bar{q}$ and VBF regions as well as the TCR, WCR, and SR. It is not true between the merged selection and resolved selection, as an event can have both a large- R jet passing a merged region definition and small- R jets passing the resolved selections. To enforce orthogonality of these two regions directly, a prioritization scheme is applied where events passing a region definition higher in the chain are not considered for those lower in the chain. The ordering of the prioritization chain is: merged SR, resolved SR, merged WCR/TCR, resolved WCR/TCR. This scheme optimizes for signal efficiency and prioritizes

the merged channel which is expected to have lower backgrounds and hence better sensitivity.

5.3 Systematic Uncertainties

This section describes the various sources of systematic uncertainty considered in the analysis. They can be divided into experimental systematics, which are evaluated from external measurements, and modeling systematics, which are evaluated specifically in the context of this analysis. In total, 92 sources of uncertainty are considered in this analysis, some of which are eigen-decompositions of a larger set of uncertainties.

The impact of each uncertainty is assessed as a function of the final discriminant used in the analysis, the invariant mass of the reconstructed diboson system $m(WV)$. The differences imposed by the uncertainties with respect to the nominal are treated as template variations to be used in the profile-likelihood approach detailed in Appendix A. The relative size of each of the systematics is controlled by a separate nuisance parameter during the fit. The profile-likelihood model is constructed in such a way that the size of conservative sources of uncertainty can be constrained to smaller values during the fit, but can not become larger than their pre-fit values. This allows the fit result itself to provide information on the sources of uncertainty which may not be estimated accurately in external measurements.

The sources of uncertainty and their estimation method are briefly described in the rest of this section. In general, an uncertainty can induce both normalization and shape effects on the predicted signal and background templates. Since the W +jets and $t\bar{t}$ normalizations are extracted from the control regions during the fit, only the shape effects of experimental systematics are considered. Due to the difficulty in concisely displaying the effects of 90+shape uncertainties, the features of the shape uncertainties are not described.

The normalization effects of the non-major backgrounds (single-top, Z +jets, SM VV , and multi-jet) and signals are characterized by Gaussian constraints with widths given in terms of percent differences with respect to the nominal prediction. A 20% normalization uncertainty then implies a Gaussian uncertainty with one standard deviations covering the range where the effected process normalization is between 0.8 or 1.2 times its nominal value.

5.3.1 Experimental Uncertainties

The majority of the experimental uncertainties are associated to the various steps of object reconstruction and calibration detailed in Section 4.3. These uncertainties are evaluated in dedicated external measurements, which are then propagated to this analysis.

Luminosity

The uncertainties on the integrated luminosity are estimated in various ways, including runs with different LHC operating conditions which are extrapolated to normal operating conditions

[84, 86]. The combined 2015+2016 luminosity uncertainty of 3.2% is applied to the background and signal normalizations estimated from MC simulation.

Pileup Reweighting

The generation of MC events typically occur before the actual data is collected. Therefore the MC is typically simulated with a $\langle\mu\rangle$ profile which does not match the data. A pile-up reweighting procedure is applied to MC to account for this effect and provide better description of data observables with MC.

Small- R Jets

The small- R jet uncertainties mainly arise from the *in-situ* jet energy scale and resolution measurements detailed in Section 4.3.2. As detailed there, several reduction schemes are provided. A 21-parameter scheme for the JES is used in this analysis. Uncertainties which are not part of the JES reduction scheme and which are treated as separate parameters include:

- Three uncertainties on the flavour composition of jets
- Three uncertainties from the η inter-calibration procedure
- Four uncertainties from the ρ pile-up removal procedure
- One uncertainty on the effects of punch-through jets
- One extrapolation uncertainty to the high- p_T regime
- One uncertainty on the JVT evaluation

Large- R Jets

The evaluation of uncertainties on the large- R jets are also detailed Section 4.3.2. The overall large- R jet p_T , mass, and $D_2^{\beta=1}$ scale calibrations are done using the R_{trk} method. The uncertainties included in the R_{trk} method are:

- Baseline uncertainty between nominal Pythia MC and data
- Uncertainty on the tracking measurement
- MC modeling uncertainty evaluated as the difference between Pythia and Sherpa dijet simulations
- Statistical uncertainty on the dijet data

Several correlation schemes are provided to reduce the number of large- R jet scale uncertainties. Each of these was tested in the final fit procedure with no noticeable differences in the results. In the end a “medium” correlation scheme was used where the p_T and mass scales are correlated, and the tracking and modeling components are uncorrelated.

b -tagging Efficiency

The evaluation of uncertainties on the b -tagging efficiencies are detailed Section 4.3.2. The variations are applied separately to the b/c /light-jets efficiency scale factors used in MC. This

analysis uses a reduction scheme where only the top three/four/five eigen-parameters are considered for the b/c /light jets. Additional uncertainties are applied for the extrapolation of the efficiency scale-factor to the high- p_T regime as well the efficiency extrapolation of c -hadron to hadronic τ decays.

Electrons and Muons

The uncertainties associated to electron and muon reconstruction are derived from tag-and-probe methods discussed in sections Section 4.3.3 and Section 4.3.4.

For both sets of objects the following uncertainties are calculated:

- Trigger efficiency scale-factor uncertainties
- Energy and momentum scale uncertainties for electrons and muons respectively
 - For muons two additional charge dependent scale uncertainties are considered
- Energy and momentum resolution uncertainties for electrons and muons respectively
 - For muons separate resolution uncertainties are evaluated for the MS and ID
- Identification and isolation uncertainties
 - For muons the following sources are treated separately: systematic, high- p_T statistical, and low- p_T statistical

For muons an additional track-to-vertex association uncertainty is applied.

Missing Transverse Energy

As discussed in Section 4.3.5, the uncertainties on the E_T^{miss} calculation are propagated from the individual object reconstruction uncertainties in the event. The effect of the soft track terms on the E_T^{miss} scale and resolution effect are considered separately and derived from $Z \rightarrow \ell\ell$ events.

An additional uncertainty is applied to events passing the E_T^{miss} trigger to account for the data and MC differences. This uncertainty includes the statistical uncertainty of the scale factor calculation as well as differences measured between $V+\text{jet}$ and $t\bar{t}$ events [215].

5.3.2 Modeling Uncertainties

Modeling uncertainties are assigned to different background and signal processes to include possible variations in the predicted normalization and shape. These systematics usually originate from uncertainty on the MC generators, such as the choice of configurations used for the generation or the choice of generator itself. These effects are quantified with dedicated studies in this analysis.

$t\bar{t}$ Modeling

The modeling uncertainty on the $t\bar{t}$ samples in the signal regions was measured by comparing the nominal MC generator with alternative generators. Since the $t\bar{t}$ normalization will be

extracted from the TCR, only the shape effects are considered. For each of the MC variations the ratio of signal region event yields, binned in $m(WV)$, with respect to the nominal generator was calculated. This ratio was then fit separately for each variation in each signal region using a linear function. This linear fit was applied as a shape systematic.

The following MC comparisons were made to calculate separate systematics:

- Matrix-element uncertainty by comparison of POWHEG with Madgraph
- Showering and hadronization uncertainty by switching Pythia with HERWIG++
- Comparison with initial and final state radiation (ISR and FSR) enriched/depleted samples where the renormalization+factorization scales and the Powheg h_{damp} parameter are adjusted. A dedicated Pythia tune is used for these samples.

V+jets Modeling

The approach of the V +jets modeling uncertainties is similar to that of the $t\bar{t}$ above. The ratio of invariant mass distributions in the signal regions of the different MC generator predictions with respect to the nominal are fit with a linear function. The list of variations considered as uncertainties are:

- Matrix-element uncertainty by a comparison with respect to Madgraph+Pythia
- Seven variations of the renormalization and factorization scale
- Two CKKW matching scale and two Resummation scale variations
- Two NNPDF30NNLO α_s variations
- 100 NNPDF30NNLO variations

Of the seven renormalization and factorization scale variations only the two with highest/lowest variation were used. The 100 NNPDF30NNLO variations are averaged using the procedure recommended by the NNPDF collaboration to provide a single uncertainty band [121].

Minor Background Uncertainties

The minor backgrounds of SM VV , single-top, and Z +jets are taken directly from MC. An overall 30% normalization uncertainty is applied to the diboson background. This is a conservative estimate of the $\sim 10\%$ expected from cross-section uncertainty and $\sim 20\%$ expected from normalization and factorization scale uncertainties. For both single-top and Z +jets, an 11% uncertainty on the normalization of is applied.

Uncertainties are derived during the fake-factor calculation to include possible variations in the multi-jet background estimation in the resolved region. An overall normalization uncertainty of 15 (30)% in the $ggF/q\bar{q}$ (VBF) region is derived from the statistical uncertainty in the fake enriched control regions.

It is worth noting that each of these background are individually less than 10% of the overall background contribution. The above quoted values then amount to percent level effects on the overall background templates.

Signal Uncertainty

Uncertainties in the MC generator calculation of the signal normalization and kinematics can effect the final excluded cross-sections extracted in Section 5.4. The PDF uncertainty can affect the normalization of signals, while uncertainties on the amount of ISR/FSR can effect the signal acceptance due to different jet multiplicities.

The PDF uncertainty on the calculated signal cross-section is evaluated by comparing the nominal choice, NNPDF3.0 [121], to MMHT2014 [123] and CT14 [124] (or CT14 to the others for the Bulk RS graviton samples). The evaluation of these uncertainties follows the procedure described in Ref. [216]. For each PDF set the 68% uncertainty band is evaluated according to each collaboration's designated procedure and the envelope of the errors is chosen as the signal uncertainty. For each signal the normalization uncertainty is derived to be less than 2%. To evaluation the ISR/FSR uncertainties, the difference between a sample with an altered Pythia [200] tune with enriched ISR/FSR contributions is used to derive normalization uncertainties of less than 3%.

5.4 Exclusion Limits on Benchmark Models

This section will illustrate the results of the resonance search in the $\ell\nu qq$ diboson channel. Binned maximum likelihood fits are done following the methods described in Appendix A, implementing all the systematics described in Section 5.3. The binned observable is the reconstructed diboson invariant mass in that region: $m(\ell\nu jj)$ in the resolved regions or $m(\ell\nu J)$ in the merged regions. The binning is optimized such that the bins are wider then the expected resonance width in HVT W' MC. The observation of a new narrow resonance should then appear as an excess of events in a single bin in each of the signal regions, or as an excess in two bins if the resonance is along a bin boundary. Bins are also merged until the expected MC statistical uncertainty in each bin is less than 70%. The binning is chosen separately for the merged and resolved regions as well as the $ggF/q\bar{q}$ and VBF regions.

There are four final fits done for each of the possible signal topologies under consideration. These explicitly are: $ggF/q\bar{q} WW$, $ggF/q\bar{q} WZ$, VBF WW , and VBF WZ . The WZ fits are interpreted in the context of the HVT W' , and the WW fits for the HVT Z' , Bulk RS graviton, and heavy Higgs-like scalar. To restate, the difference between WW and WZ selection is whether it passes the respective V -tagger selection in the merged channel, or the the respective dijet mass window in the resolved channel. In each of the fits the merged HP, merged LP, and resolved signals regions and control regions are fit simultaneously. Common systematics between the regions are correlated. The normalization for each process is treated as correlated between all regions, with the exception of the W +jets and $t\bar{t}$ which are separated between the merged and resolved regions.

Figures 5.5 - 5.6 show the data and fitted background MC distributions in all signal regions. Good agreement between data and background is found in all signal regions. The smallest local

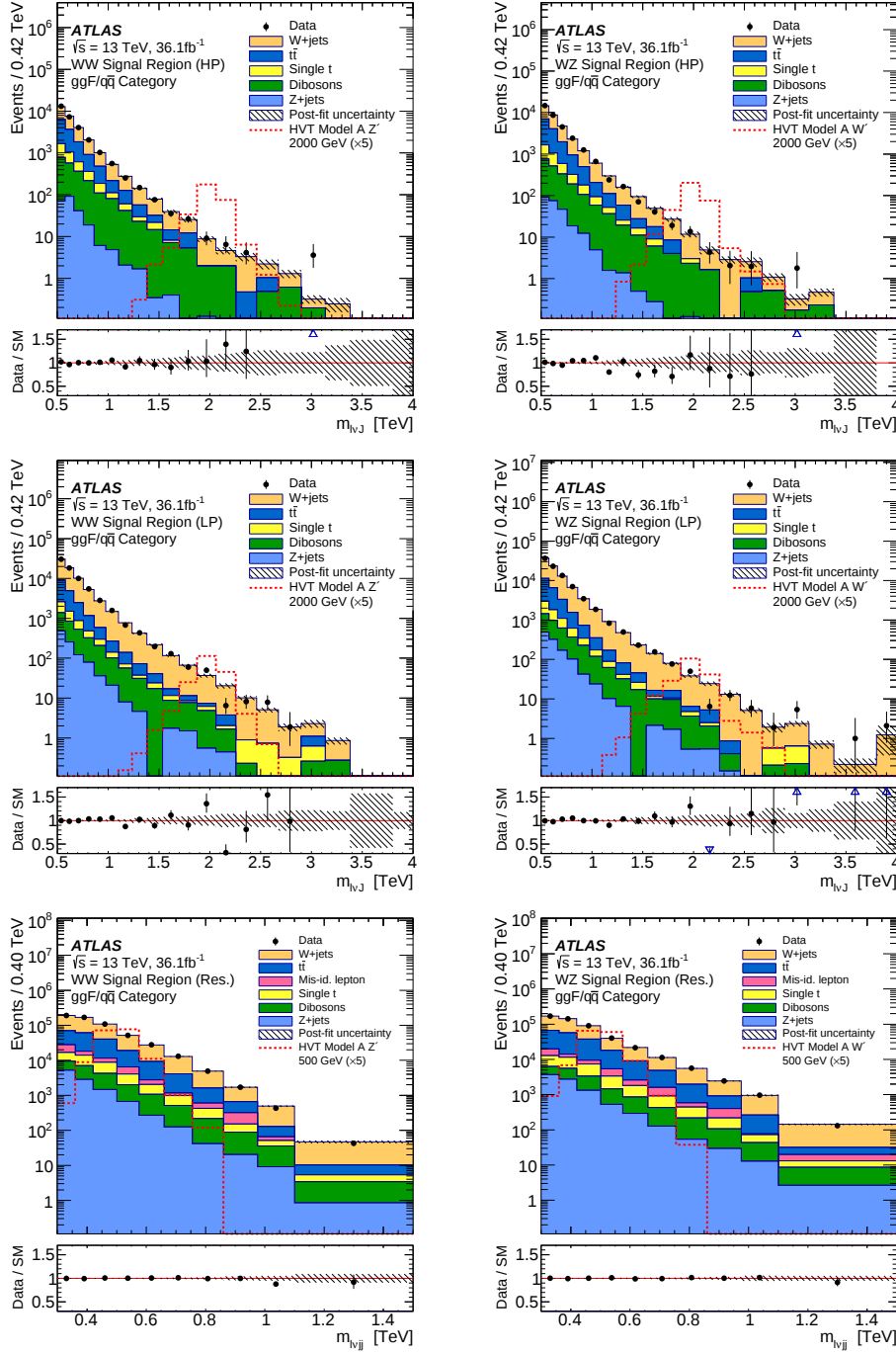


Figure 5.5: The fitted $m(\ell\nu jj)$ or $m(\ell\nu J)$ distributions for merged HP, merged LP, and resolved (Top, Middle, Bottom), WW and WZ (Left, Right) $ggF/q\bar{q}$ signal regions. The data is shown in black and the filled regions show the separated contributions of the measured background processes. The total uncertainty, including both systematics and statistical uncertainties is shown as black hashed region on the background envelope. The distribution of a resonance as predicted for a HVT Model A Z' of mass 2 TeV (or 0.5 TeV in the resolved plots) is shown with the cross-section scales by a factor of 5 for increased visibility. The bottom plots display the ratio of SM background prediction to data in that bin. The binning shown was used for the final hypothesis testing.

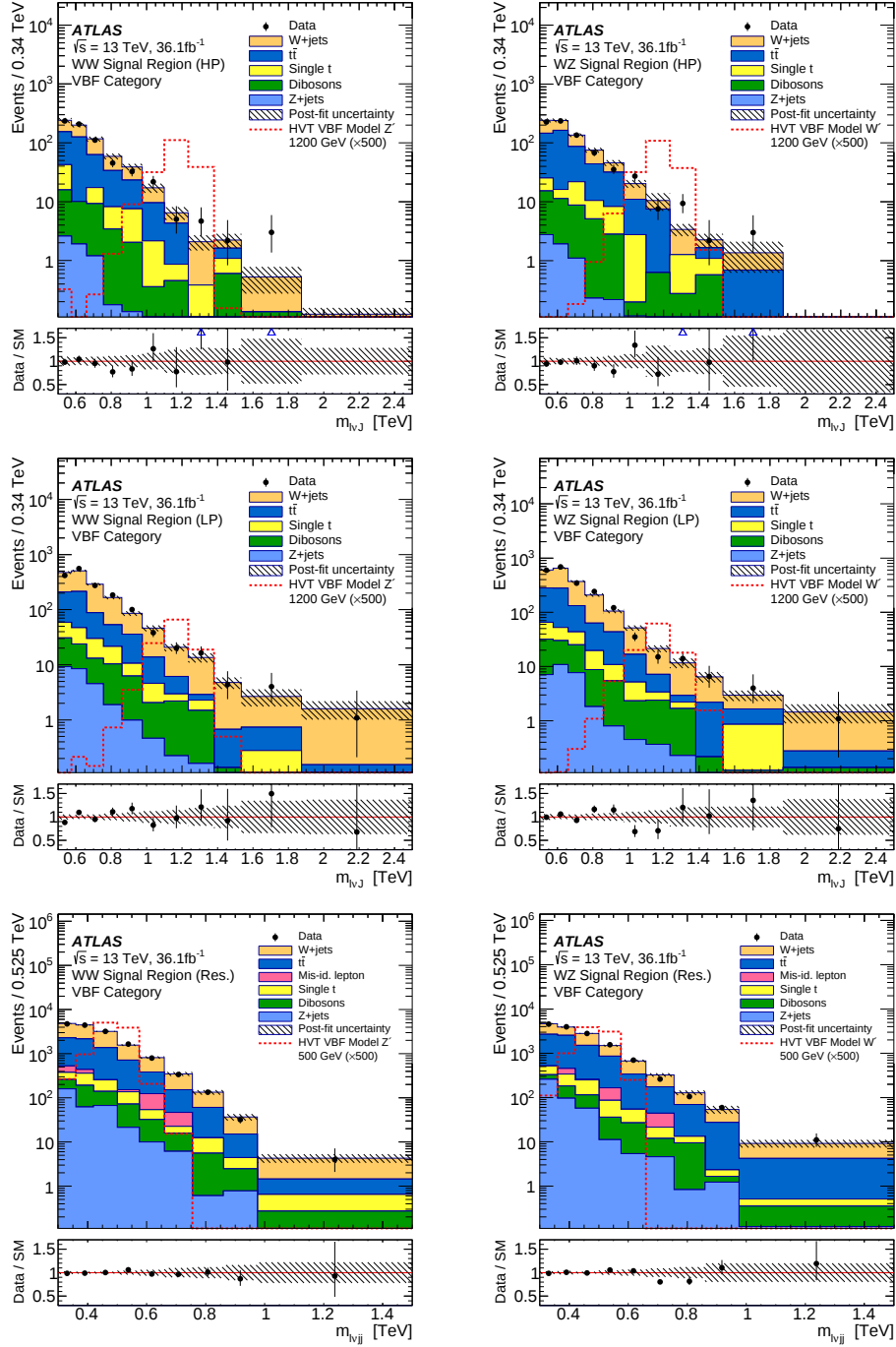


Figure 5.6: The fitted $m(\ell\nu jj)$ or $m(\ell\nu J)$ distributions for merged HP, merged LP, and resolved (Top, Middle, Bottom), WW and WZ (Left, Right) VBF signal regions. The data is shown in black and the filled regions show the separated contributions of the measured background processes. The total uncertainty, including both systematics and statistical uncertainties is shown as black hashed region on the background envelope. The distribution of a resonance as predicted for a HVT Model A Z' of mass 2 TeV (or 0.5 TeV in the resolved plots) is shown with the cross-section scales by a factor of 5 for increased visibility. The bottom plots display the ratio of SM background prediction to data in that bin. The binning shown was used for the final hypothesis testing.

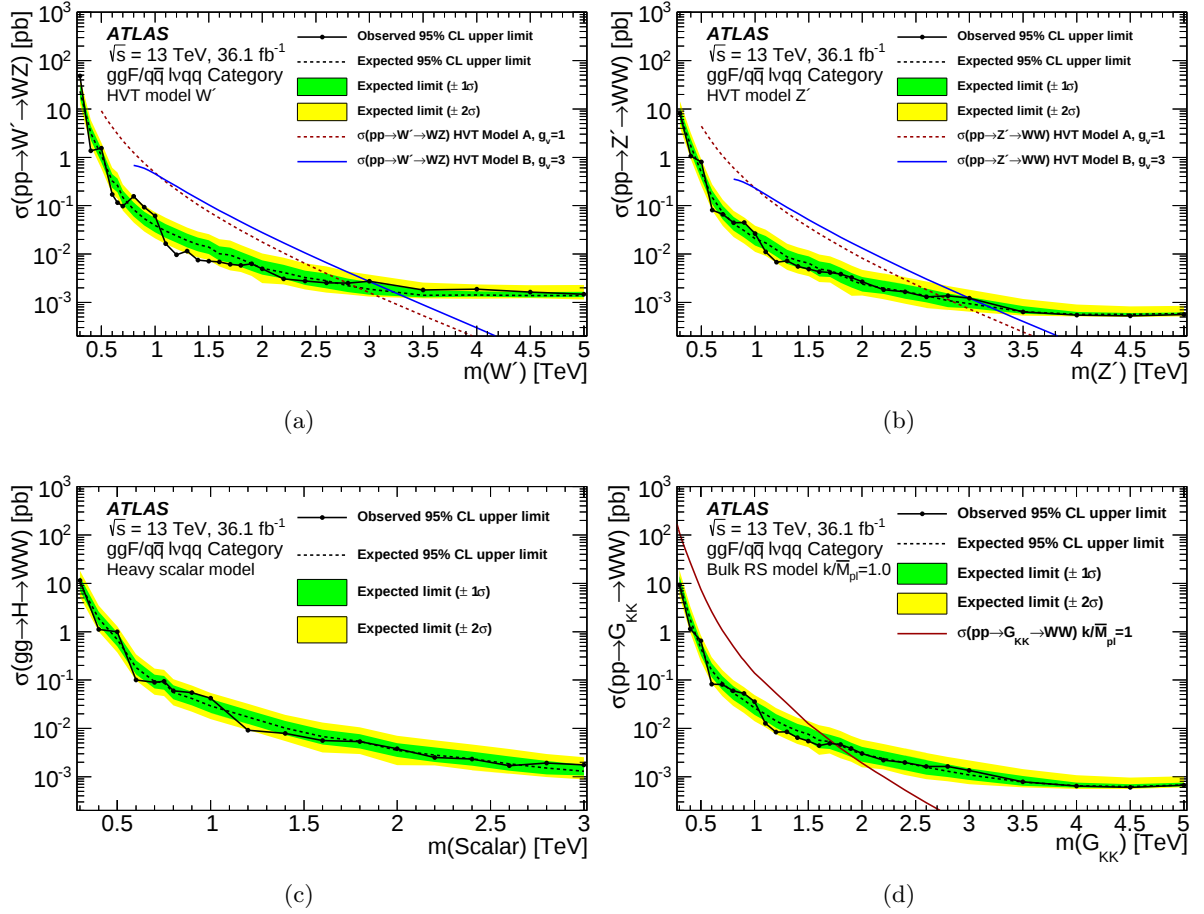


Figure 5.7: The 95% CL_s upper limits on the cross-section times branching ratio as a function of resonant mass for the a) HVT W' , b) HVT Z' , c) heavy Higgs-like scalar, and d) Bulk RS graviton as produced either by ggF or $q\bar{q}$ annihilation. The solid black curves shows the observed limit, and the dashed black curve shows the expected limit, with one and two sigma confidence bands in green and yellow. Theoretical cross-sections for relevant models are shown in blue or red.

p -value for the background only hypothesis corresponds to a significance² of 2.7σ in the WW VBF regions (left column of Figure 5.6) at 1.6 TeV. Additional studies on the fit model used in this analysis and local background only p -value calculations can be found in Appendix B.

Exclusion limits on the signal cross-section in the signal+background hypothesis are calculated for each benchmark model interpretation using the CL_s method at 95% confidence interval. The test statistic used is given in Equation A.6 and is evaluated separately for several mass points in each model. The excluded upper limit on the cross-sections multiplied by the branching ratio for the considered signal models are shown in Figures 5.7 - 5.8 for different signal masses. The ggF/ $q\bar{q}$ HVT and Bulk RS graviton interpretations have theoretical predictions overlaid. The heavy Higgs-like scalar interpretation is not specific to a single model and

²The relation between p -values and significance values are detailed in Appendix A.1.

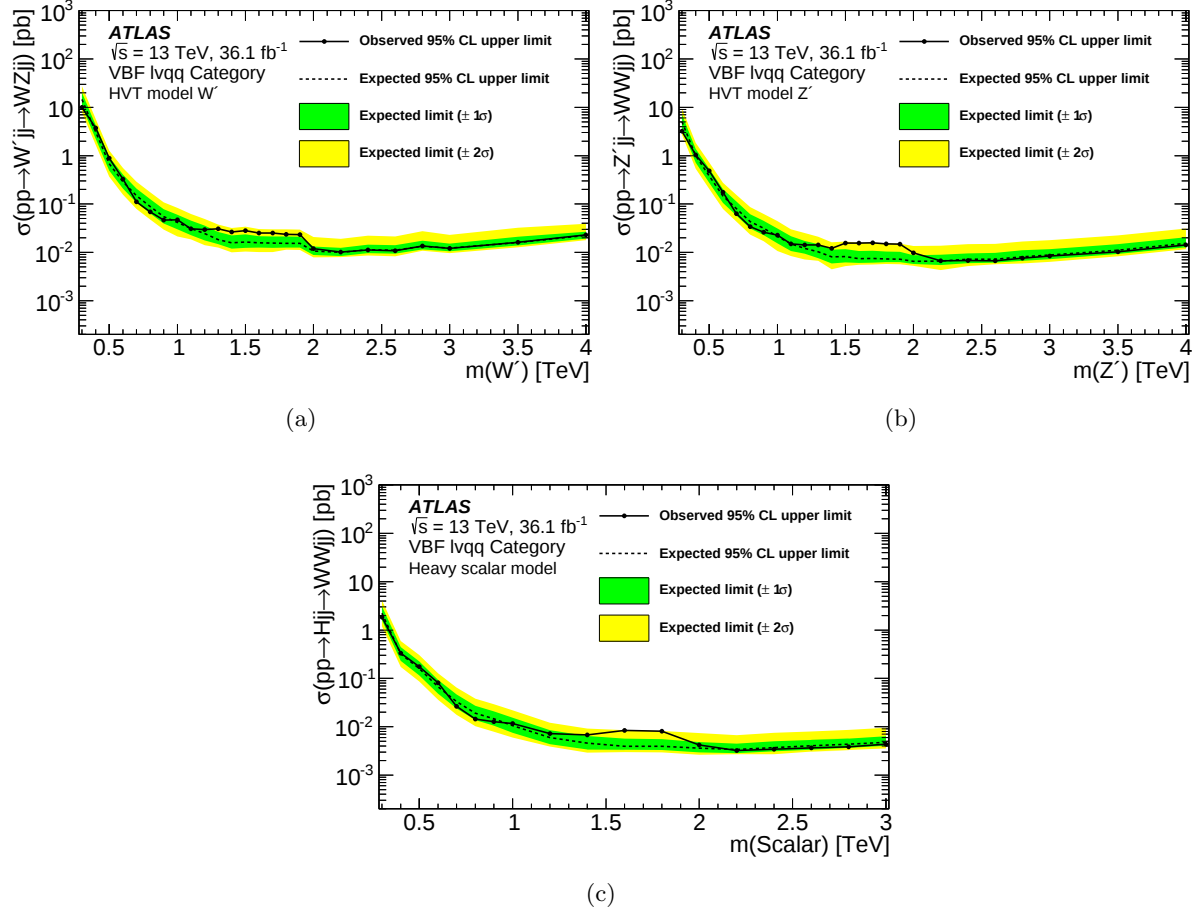


Figure 5.8: The 95% CL_s upper limits on the cross-section times branching ratio as a function of resonant mass for the a) HVT W' , b) HVT Z' , and c) heavy Higgs-like scalar as produced via VBF. The solid black curves shows the observed limit, and the dashed black curve shows the expected limit, with one and two sigma confidence bands in green and yellow. Theoretical cross-sections for relevant models are shown in blue or red.

VBF Category			
$m(Z') = 1200$ GeV		$m(W') = 500$ GeV	
Source	$\Delta\mu/\mu$ [%]	Source	$\Delta\mu/\mu$ [%]
MC statistical uncertainty	15	MC statistical uncertainty	16
Large- R jets mass resolution	5	W +jets: cross section	10
W +jets: PDF choice	5	Multijet E_T^{miss} modeling	10
$t\bar{t}$: alternative generator	5	Small- R jets energy resolution	9
W +jets: cross section	5	SM diboson cross section	8
$t\bar{t}$: scales	4	$t\bar{t}$: cross section	7
Total systematic uncertainty	24	Total systematic uncertainty	40
Statistical uncertainty	52	Statistical uncertainty	30

ggF/ $q\bar{q}$ Category			
$m(W') = 2000$ GeV		$m(Z') = 500$ GeV	
Source	$\Delta\mu/\mu$ [%]	Source	$\Delta\mu/\mu$ [%]
MC statistical uncertainty	12	Large- R jets kinematics	17
W +jets: generator choice	8	MC statistical uncertainty	12
W +jets: scale	5	$t\bar{t}$: scale	11
SM diboson normalization	4	SM diboson cross section	10
Large- R jets mass resolution	4	W +jets: alternative generator	10
Large- R jets D_2 resolution	4	W +jets: scale	9
Total systematic uncertainty	20	Total systematic uncertainty	42
Statistical uncertainty	50	Statistical uncertainty	18

Table 5.4: The relative impact of the six topmost ranking sources of uncertainty for two mass points in the VBF channel (top) and ggF/ $q\bar{q}$ channel (bottom). The relative impact is measured by varying a single probed systematic by $\pm 1\sigma$ with respect to the nominal prediction and measuring the change on the signal strength $\Delta\mu$ before and after the change.

can be interpreted in the context of Two-Higgs Doublet Model of Section 2.3.1 or the radion of RS models in Section 2.3.3. The data excludes at the 95% CL_s confidence level new HVT Z' resonances with masses below 2.7 (3.0) and HVT W' resonances below 2.8 (3.0) TeV in the HVT Model A (B). For the Bulk RS graviton, the upper mass limit is found to be 1.7 TeV for coupling $k/M_{Pl} = 1.0$.

The breakdown of the top sources of uncertainty can be seen in Table 5.4 for several mass points of both ggF/ $q\bar{q}$ and VBF interpretations. The largest sources of uncertainty for low mass points at 500 GeV originate from MC modeling and jet based experimental systematics. At higher masses of 1 - 2 TeV the dominant systematic uncertainty is the limited MC generator statistics for the background samples and large- R jet uncertainties. For masses above 1 TeV, the statistical uncertainty on data is the limiting factor.

Summary

This section summarizes the results obtained from a search of new narrow resonances decaying to diboson pairs with 36 fb^{-1} of ATLAS data [8]. In comparison to previous searches in the $\ell\nu q\bar{q}$ channel [187, 188], this includes a full harmonization of techniques for reconstructing the

$V \rightarrow q\bar{q}$ decays. This publication also included a new interpretation of the search for VBF produced resonances. The preceding publication from ATLAS on diboson resonance results used 3.6 fb^{-1} of 13 TeV data, but utilized four decay channels: $VV \rightarrow \ell\nu qq + \ell\ell qq + \nu\nu qq + qqqq$ [189]. The results shown in this chapter only used one of these channels ($\ell\nu qq$), but improved on the excluded upper mass limits of the HVT resonances by 100 GeV.

The main results and analysis procedure of this work form the backbone of several different diboson studies discussed in subsequent chapters. Chapter 6 combines the displayed $\ell\nu qq$ results with all concurrent diboson and dilepton resonance searches to provide the strongest constraints on new resonant physics models from ATLAS with 36 fb^{-1} of 13 TeV data. The overall analysis procedure used for the new VBF interpretation is adapted in Chapter 7 to provide a measurement on the cross-section of electroweak vector-boson scattering in the semi-leptonic channels. Lastly, Chapter 8 presents prospect studies of this analysis and those of Chapter 7 in the HL-LHC era.

Chapter 6

Diboson+Dilepton Resonance Combination with 36 fb^{-1}

Alone we can do so little; together we can do so much.

—Hellen Keller

Publications for resonance searches typically target a specific final state, as was done in Chapter 5 for the diboson search in the $\ell\nu qq$ channel. This is due to different final states requiring different experimental techniques and dedicated teams of researchers working on each. In general though, one does not usually expect a new resonance to occur in a specific channel only. For example, evidence of WZ resonances should appear in all relevant final states, as they differ only by SM W/Z branching ratios. It is also true that general new physics does not need to solely couple to diboson states, and can be evident in other decay modes, such as to pairs of leptons or quarks. Each of the individual analyses are then typically providing complementary information on similar models, which is not exploited in the individual publications.

It is useful to statistically combine several search analyses for multiple reasons. The main reason is that the combination will provide equal or better signal sensitivity than any of the analyses individually. For example, if multiple analyses have comparable sensitivity, then the combination will provide a signal sensitivity approximately proportional to the individual analyses values summed in quadrature. If multiple analyses have varying sensitivity over regions of phase-space, then the combined result will provide the most stringent constraints over the whole phase-space. Another useful result provided by combining analyses is that it can check whether multiple excesses or deficits coincide. Individual analyses may see such features, but not be able to claim any strong result due to limited statistics. The combination may be able to claim evidence for such features, if they align in several individual analyses. Conversely, bumps and excesses can be smoothed in the combination, which can provide the insight that these features were most likely statistical fluctuations in the individual analyses.

This section will detail the statistical combination of existing ATLAS searches with 36 fb^{-1} in diboson and dilepton final states [9]. In this context, diboson will refer to states with pairs of

electroweak gauge bosons (VV), or one W/Z boson and a Higgs boson (Vh). Dilepton searches will refer to those searching for pairs of opposite-sign same flavour electrons and muons ($\ell\ell$), or one electron or muon and $E_{\text{T}}^{\text{miss}}$ originating from a neutrino ($\ell\nu$)¹. Previously, both the ATLAS [189] and CMS [217] collaborations presented similar results with under 3 fb^{-1} of 13 TeV data, but only in the context of diboson channel combinations. The analysis presented here was the first from either collaboration to simultaneously combine diboson with dilepton channels. After this publication, the CMS collaboration published a similar combination of diboson and dilepton results [218].

This analysis combines 14 final states which are separated over 10 individual publications. Resonant Vh searches have only been conducted for the channels with the SM Higgs decaying as $h \rightarrow b\bar{b}$, while all resonant VV channels were explored. The total list of final states includes:

- $X \rightarrow VV \rightarrow \ell\nu qq$ [8]
- $X \rightarrow VV \rightarrow \ell\ell qq + \nu\nu qq$ [191]
- $X \rightarrow VV \rightarrow qq qq$ [192]
- $X \rightarrow VV \rightarrow \ell\nu\ell\nu$ [194]
- $X \rightarrow VV \rightarrow \ell\nu\ell\ell$ [193]
- $X \rightarrow VV \rightarrow \ell\ell\ell\ell + \ell\ell\nu\nu$ [195]
- $X \rightarrow Vh \rightarrow \ell\ell b\bar{b} + \ell\nu b\bar{b} + \nu\nu b\bar{b}$ [219]
- $X \rightarrow Vh \rightarrow qq b\bar{b}$ [220]
- $X \rightarrow \ell\ell$ [221]
- $X \rightarrow \ell\nu$ [222]

Each of the contributing analyses were reviewed and published individually. We assume that each of the analyses appropriately estimated all background contributions and evaluated the relevant uncertainties. The individual analysis results and fit models are used “as-is”, and any changes made will be explicitly detailed. In the rest of this chapter only the relevant details of the statistical combination will be discussed. The overall approach is to combine the individual profile-likelihood fits discussed in Appendix A by simultaneously fitting all the analyses accounting for correlations between relevant systematics. The signal models used for the interpretation of data are discussed in Section 6.1. Relevant details for the combination in regards to event orthogonality and treatment of correlations are discussed in Sections 6.2 - 6.3. Finally, the combined results will be interpreted for the benchmark models in Section 6.4 and directly on the coupling parameters of the HVT model in Section 6.5.

6.1 Signal Models

The combination targets the same set of models as considered in the diboson resonance searches detailed in Chapter 5. Explicitly, the 3 models considered are:

- Spin-0: Heavy Higgs-like scalar H with narrow width

¹Anti-particle labels are suppressed in the labeling of analyses.

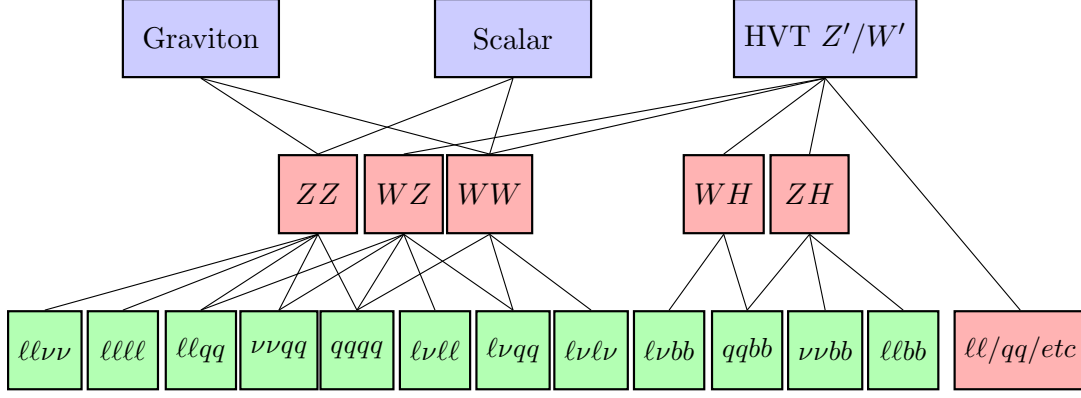


Figure 6.1: Web-graph of the signal models (blue) considered connected to the decay modes (red) and possible tertiary states (green).

- Spin-1: Heavy Vector Triplet (HVT) W'/Z'
- Spin-2: Bulk Randall-Sundrum (RS) graviton G_{KK}

The heavy Higgs-like scalar and Bulk RS graviton models are only used for interpretation of the $WW + ZZ$ type resonances. Almost all final states are used for the HVT $W' + Z'$ combination, except for those mediated by ZZ decays as there is no tree-level ZZ resonances in the HVT model. A sketch of which signal models are considered in which final state is shown in Figure 6.1. The HVT signals are interpreted in the three benchmark models designated A, B, and C, which are detailed in Section 2.3.2. Models A and B will be used to probe $q\bar{q}$ production modes and Model C will be used for the VBF production. The Bulk RS graviton is interpreted with the model parameter $k/M_{Pl} = 1.0$.

For all models, the resonances are assumed to be sufficiently narrow such that the effect of interference between signal and backgrounds are negligible. The sensitivity to interference is channel and parameter dependent. For the heavy Higgs-like scalar decaying in the leptonic channel, the interference can change the cross-section by approximately 10% for widths of $\Gamma/m < 0.15$ [223]. The combination targets the parameter space with signals widths $\Gamma/m < 0.05$ where the effects of interference on signal normalizations and shape is negligible.

Similar to Chapter 5, all models are interpreted in the context of gluon-gluon fusion (ggF) or quark-antiquark ($q\bar{q}$) production. Signals produced in the vector-boson fusion (VBF) production mode are also interpreted for the heavy Higgs-like scalar and HVT model, but only in the VV channels. No attempt at combined ggF/ $q\bar{q}$ + VBF interpretations is investigated as the individual analyses applied non-complementary approaches to fits of the VBF signal. Several analyses applied the most robust approach and fit both signal production types simultaneously in the ggF/ $q\bar{q}$ and VBF channels with adjustable normalizations, while others only conducted dedicated fits in each channel. Thus, there is not sufficient information from each of the analyses to do a global ggF/ $q\bar{q}$ + VBF interpretation. Each channels approach is summarized in Table 6.1. Separate ggF/ $q\bar{q}$ and VBF results will be shown for the relevant models.

It is also worth noting that not all the individual analyses covered the same parameter

Input channel	Mass Range [GeV]	Type of background fit	VBF Treatment
$WW/ZZ/WZ \rightarrow qq\bar{q}\bar{q}$	1200-5000	fit with a parametric function	no VBF signal
$ZZ/WZ \rightarrow \nu\nu q\bar{q}$	500-5000	binned fit using MC templates	fit $ggF/q\bar{q}$ +VBF signal in both regions
$WW/WZ \rightarrow \ell\nu q\bar{q}$	500-5000	binned fit using MC templates	fit $ggF/q\bar{q}$ (VBF) signal in $ggF/q\bar{q}$ (VBF) region
$ZZ/WZ \rightarrow \ell\ell q\bar{q}$	500-5000	binned fit using MC templates	fit $ggF/q\bar{q}$ +VBF signal in both regions
$ZZ \rightarrow \ell\ell\nu\nu$	300-2000	binned fit using MC templates	fit both $ggF/q\bar{q}$ and VBF separately in both regions
$WW \rightarrow \ell\nu\ell\nu$	300-3000	binned fit using MC templates	fit both $ggF/q\bar{q}$ and VBF separately in both regions
$WZ \rightarrow \ell\nu\ell\ell$	300-3000	binned fit using MC templates	fit $ggF/q\bar{q}$ (VBF) signal in $ggF/q\bar{q}$ (VBF) region
$ZZ \rightarrow \ell\ell\ell\ell$	300-1200	fit with a parametric function	fit both $ggF/q\bar{q}$ and VBF separately in both regions
$WH/ZH \rightarrow qq\bar{b}\bar{b}$	1200-5000	fit with a parametric function	no VBF signal
$ZH \rightarrow \nu\nu b\bar{b}$	500-5000	binned fit using MC templates	no VBF signal
$WH \rightarrow \ell\nu b\bar{b}$	500-5000	binned fit using MC templates	no VBF signal
$ZH \rightarrow \ell\ell b\bar{b}$	500-5000	binned fit using MC templates	no VBF signal
$\ell\nu$	200-5500	binned fit using MC templates	no VBF signal
$\ell\ell$	200-5500	binned fit using MC templates	no VBF signal

Table 6.1: Summary of different approaches used in the analyses. The explored mass range, approach to background modeling, and treatment of VBF signals are shown.

space. Individual analyses only published results for the resonant mass range they are most sensitive to as seen in Table 6.1. We will find in Section 6.4, that the exclusion limits on the HVT Model A benchmark will extend beyond the 5 TeV point where the diboson resonances searches stopped. To extend the diboson limits to the same mass range as the dilepton analyses, the total $VV+Vh$ combination is extrapolated up to 5.5 TeV by rescaling the 5 TeV results by the expected differences in the efficiency and cross-section. This extrapolation method is justified as above 5 TeV the diboson analyses have no expected or observed background events, and limits are dominated by expected signal yields.

6.2 Orthogonality

The statistical combination can be greatly simplified if all contributing analyses have orthogonal selection with respect to each other. If this is true, no single data event can be selected by multiple analyses. There is then no correlation of data events between channels, and they can be considered as separate independent categories in the fit. If this is not true, correlations must be evaluated and applied, which is a time-consuming and practically difficult task. The aim of this section is to detail the studies which went into understanding and handling overlaps between different signal regions. Orthogonality of control regions is not imposed in the combination.

Due to the scope of this combination, the cuts imposed in each individual analysis will not

Channel	Diboson state	Selection				VBF cat.
		Leptons	E_T^{miss}	Jets	b -tags	
$qqqq$	$WW/WZ/ZZ$	0	veto	$2J$	—	—
$\nu\nu qq$	WZ/ZZ	0	yes	$1J$	—	yes
$\ell\nu qq$	WW/WZ	$1e, 1\mu$	yes	$2j, 1J$	—	yes
$\ell\ell qq$	WZ/ZZ	$2e, 2\mu$	—	$2j, 1J$	—	yes
$\ell\ell\nu\nu$	ZZ	$2e, 2\mu$	yes	—	0	yes
$\ell\nu\ell\nu$	WW	$1e+1\mu$	yes	—	0	yes
$\ell\nu\ell\ell$	WZ	$3e, 2e+1\mu, 1e+2\mu, 3\mu$	yes	—	0	yes
$\ell\ell\ell\ell$	ZZ	$4e, 2e+2\mu, 4\mu$	—	—	—	yes
$qqbb$	WH/ZH	0	veto	$2J$	1, 2	—
$\nu\nu bb$	ZH	0	yes	$2j, 1J$	1, 2	—
$\ell\nu bb$	WH	$1e, 1\mu$	yes	$2j, 1J$	1, 2	—
$\ell\ell bb$	ZH	$2e, 2\mu$	veto	$2j, 1J$	1, 2	—
$\ell\nu$	—	$1e, 1\mu$	yes	—	—	—
$\ell\ell$	—	$2e, 2\mu$	—	—	—	—

Table 6.2: Summary of analysis channels, possible intermediate boson state, and experimental signatures. The notation “ j ” represents small- R jets and “ J ” large- R jets. Leptons are either electrons “ e ” or muons “ μ ”. Categories with — indicate such requirements are not imposed. Additional jets (not included in the “Jets” column) are required to define VBF categories.

be discussed explicitly. The exact details are found in each of the individual publications. A general overview of the objects selected in each of the analyses can be seen in Table 6.2.

All of the VV analyses are by design orthogonal to each other and the same is true for the Vh channels with respect to each other. The VV and Vh channels are not mutually orthogonal though, and this overlap is considered in Section 6.2.1. The dilepton channel is also not by design orthogonal to the VV/Vh channels and the overlap is discussed in Section 6.2.2.

6.2.1 VV and Vh Overlap

The VV and Vh analysis which have hadronically decaying W/Z and Higgs bosons are not orthogonal due to limited detector resolution for reconstructed jets and the imperfect identification of b -tagged jets. This induces significant overlap between analyses such as $\ell\ell qq$ and $\ell\ell bb$.

Both sets of analyses reconstruct the $W/Z \rightarrow q\bar{q}$ and $h \rightarrow b\bar{b}$ decays as either two small- R jets or one large- R jet. These are called the resolved and merged selection respectively. For the Vh analyses at least one of the small- R jets is b -tagged and one track-jet associated to the large- R jet is b -tagged. The VV analysis also allows b -tagged jets for the resolved channels to reconstruct the $Z \rightarrow b\bar{b}$ decays, and no explicit b -tagging is done in the merged channels.

The overlap between these channels is only at the object selection level and could be mitigated if subsequent analysis cuts enforced orthogonality. For the VV and Vh analyses this is not true for two reasons. The first is the invariant mass window cuts applied on the dijet system and large- R jet overlap between the two analysis. The mass windows applied in each of the

	$\ell\ell bb$	$\ell\nu bb$	$\nu\nu bb$	$qqbb$
Resolved	[100,145]	[110,140]	[110,140]	
Merged	[75,145]	[75,145]	[75,145]	[75,145] for H [70,110] for Z
	$\ell\ell qq$	$\ell\nu qq$	$\nu\nu qq$	$qqqq$
Resolved	[70,105] for Z [62,97] for W	[82,106] for Z [66,94] for W		
Merged	[76,106] for Z [65,95] for W	[69,114] for Z [64,104] for W	[76,106] for Z [65,95] for W	[76,106] for Z [65,95] for W

Table 6.3: Hadronic W/Z and Higgs mass windows used for the Vh (top) and VV (bottom) analysis. The Z/W mass windows for large- R jets in the boosted channels are p_T dependent so the largest window is quoted.

relevant analysis are shown in Table 6.3. The VV resolved channels and all Vh channels apply a fixed window cut, while for the VV merged analysis a p_T dependent window cut is applied. The second source of overlap is the individual analysis choice of internal overlap removal between the merged and resolved selections, which differs between the VV and Vh channels. To handle the situation where an event contains both two resolved jets and one large- R jet satisfying all the selection criteria, both channels apply a prioritization scheme which removes overlapping events in all but one of the signal/control regions. Such a scheme was discussed in context of the $\ell\nu qq$ search in Section 5.2.6. The VV analyses prioritize the merged channel, while the Vh prioritizes the resolved.

Orthogonality can be re-introduced by applying some new prioritization scheme as was done internally within each analysis (as discussed in Section 5.2.6 for the $\ell\nu qq$ search for example). This entails a choice of either removing events which pass the VV selection from Vh or vice-versa. Since the VV analyses are slightly more sensitive, they are prioritized. Any event which passes a VV analysis is therefore removed from consideration of the Vh analyses. Figure 6.2 shows the change in sensitivity between the default analysis and the updated analysis with the vetoed events. The differences can be seen to be minimal.

An important feature of this choice of orthogonality scheme, is that signal events are not lost by introducing the veto. Instead signal events which would have fallen in the overlapping region of the Vh analyses migrate to the VV signal region. In the subsequent combination of $VV+Vh$ the overall sensitivity is recovered.

When displaying results below, if the Vh channels are solely shown, these will refer to a combination of the default analyses. When any combination of Vh channels with other channels is shown, the altered result with vetoed events will be used.

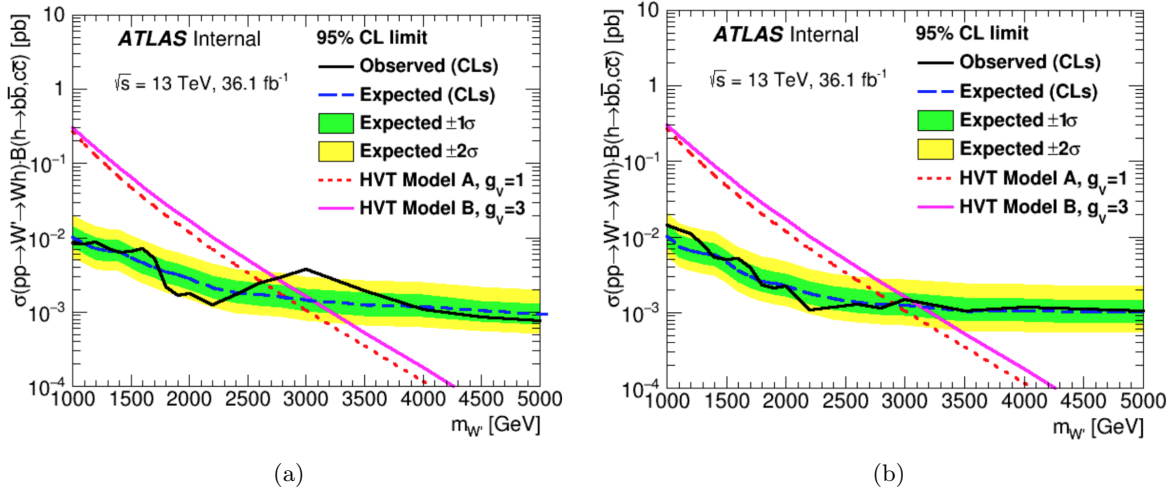


Figure 6.2: The HVT limits provided by the default Vh analysis (left) and the Vh analysis with the overlapping VV events removed (right).

6.2.2 Dilepton and Diboson Overlap

An overlap in events is also found between the diboson and dilepton analyses. This originates from the fact that the dilepton events are inclusive with respect to additional objects in the event. For example, no veto on additional jets is applied in the $\ell\nu$ analysis to remove overlap with the $\ell\nu q\bar{q}$ events.

The number of overlapping data events between the signal regions of the VV and dilepton analyses can be seen in Figure 6.3. With respect to the total number of events considered in each analysis, the number of overlapping events is small. Additional checks were made as the effect of this overlap could still potentially be non-negligible if the overlapping events are found in the low statistics tails of the distributions. In the dilepton channels it is found that the $O(10)$ overlapping events only impact the low-mass regions where $O(10^4)$ background events are found. The overlap is therefore a sub-percent level effect. In contrast to the VV and Vh overlaps detailed in Section 6.2.1, these overlaps can be safely ignored. Technically, this can be considered as a prioritization scheme where the VV is prioritized over the dilepton channels, but since the effect is small, it can be effectively ignored in the dilepton channels.

6.3 Correlation Scheme

The majority of the systematic uncertainties in the individual analyses are evaluated in the same way; from external measurements during object reconstruction and calibration studies. The systematics are then expected to be correlated amongst the individual analyses. Ambiguity can arise if a common source of uncertainty is estimated through different means, or if different nuisance parameter reduction schemes are used for the small- R jet systematics for example.

	$qqqq$	$\ell\ell qq$	$\ell\nu qq$	$\nu\nu qq$	$\ell\nu\ell\nu$	$\ell\ell$	$\ell\nu$
$qqqq$	1304						
$\ell\ell qq$		23336					25
$\ell\nu qq$			14419				4
$\nu\nu qq$				10373			
$\ell\nu\ell\nu$					11305		
$\ell\ell$						115819	
$\ell\nu$		25	4				136501

(a)

	$qqqq$	$\ell\ell qq$	$\ell\nu qq$	$\nu\nu qq$	$\ell\ell\ell$	$\ell\ell\nu$	$\ell\nu\ell\nu$	$\ell\ell$	$\ell\nu$
$qqqq$	1180								
$\ell\ell qq$		22175							28
$\ell\nu qq$			11929						3
$\nu\nu qq$				3139					
$\ell\ell\ell$					560		1	24	
$\ell\ell\nu$						292			
$\ell\nu\ell\nu$					1		11305		
$\ell\ell$					24			115819	
$\ell\nu$		28	3						136501

(b)

Figure 6.3: Number of data events overlapping between the signal regions of the respective channels. The diagonal terms show the total data yields in that channel. The $WW + WZ$ plot (left) is relevant for the HVT interpretation while the $WW + ZZ$ (right) is relevant for the heavy Higgs-like scalar and Bulk RS graviton interpretation.

To apply a consistent approach throughout the whole combination, a correlation scheme is used where any systematic which is estimated from the exact same source and method is 100% correlated. If two systematics represent the same source of uncertainty but are estimated differently, they are completely uncorrelated.

None of the individual analyses considered sources of experimental uncertainties which have not been discussed previously in Section 5.3 for the $\ell\nu qq$ analysis, but they may differ in their evaluation method or eigenvalue reduction scheme. Additional modeling uncertainties are also commonly considered in non- $\ell\nu qq$ analyses. A full list of systematics in each analysis and their differences with respect to the $\ell\nu qq$ analysis are detailed in in Appendix B.3.

To summarize the correlation scheme for the object systematics, all electron and muon systematics are correlated amongst all channels, except for the $\ell\nu\ell\nu$ channel which has two electron identification components. Similarly, all channels with large- R jets use the same “medium” reduction scheme and are correlated. In regards to the small- R jet uncertainties, most channels use the 21 parameter reduction scheme which are correlated. The $\ell\nu\ell\ell$, $\ell\ell\ell\ell$, and $\ell\ell\nu\nu$ use the 3 nuisance parameter scheme and are correlated between each other but uncorrelated from the rest. Both small- R jets and track-jets are used for b -tagging, with the associated uncertainty on the tagging considered uncorrelated between the collections.

Modeling systematics are treated in the same way. The estimated signal modeling is correlated between the hadronic and semi-leptonic final states since the ISR/FSR contribution and the PDF variation are estimated in the same way. The leptonic channels include various addi-

tional theoretical uncertainties which are uncorrelated. The $t\bar{t}$ shape modeling are consistently estimated in the $\ell\ell qq$, $\ell\nu qq$, and $\nu\nu qq$ channels and are correlated. The normalization of the $t\bar{t}$ background are uncorrelated between channels. Similarly for the V +jets background, the $\ell\nu qq$ and $\nu\nu qq$ use PDF/scale variations while the $\ell\ell qq$ uses a data-driven estimate, hence they are uncorrelated. Diboson systematics are estimated differently in the leptonic and semi-leptonic channels and are uncorrelated. Lastly fake-lepton estimations are fully uncorrelated between relevant analyses.

6.4 Exclusion Limits on Benchmark Model

In this section, the results of the combination of resonance searches are presented for the benchmark models discussed in Section 6.1. All models are fit using a likelihood model with additional nuisance terms parameterizing the uncertainties, as detailed in Appendix A. The 95% CL_s upper limits on the cross-section for each signal model and mass in the signal+background hypothesis are evaluated. Additional studies on the fit models used in the combination and local background only p -value calculations can be found in Appendix B.

The excluded 95% CL_s upper limits for WW , WZ , and ZZ combinations are shown in Figures 6.4 - 6.8 for different signal interpretations. The breakdown of the expected limits from the individual channels with respect to the total combination are also shown. Theory curves are overlaid for the Bulk RS graviton interpretations as well as the HVT Models A and B for the $q\bar{q}$ interpretations and Model C for the VBF interpretations. It is found that for resonant masses larger than 1 TeV, the semileptonic and fully hadronic channels provide the strongest constraints, while for resonant masses on the scale of hundreds of GeV, the leptonic channels dominate. As expected, the combined limits provide the strongest limits over the whole mass range, with the improvement proportional to the number of channels in each fit and their relative sensitivity. It can be seen that for some fits, such as the HVT W' interpretation in Figure 6.4b, the 95% CL_s excluded cross-section times branching ratio at the mass point of 5 TeV is improved by a factor of seven in the combination. For other combinations where one channel dominates over the whole mass range, such as the VBF heavy Higgs-like scalar interpretation in Figure 6.7a, no improvement is found.

The upper limits for the combined VV fits in all interpretations can be seen in Figure 6.9. In comparison with the results of Section 5.4, the observed upper mass limit at 95% CL_s for the $q\bar{q}$ produced HVT W' signal in Model A (B) increases from 2.8 (3.2) TeV in the $\ell\nu qq$ channel to 3.7 (4) TeV in the VV combination. Similarly, the observed upper limits for the Bulk RS graviton increase from 1.6 TeV in the $\ell\nu qq$ channel to 2.3 TeV in the combination. The observed increase is not as sizable for the graviton model as the theoretically expected cross-section decreases very sharply as a function of mass in comparison to the HVT model prediction.

The observed and excluded limits in the HVT model from the Vh , dilepton, and total dibo-

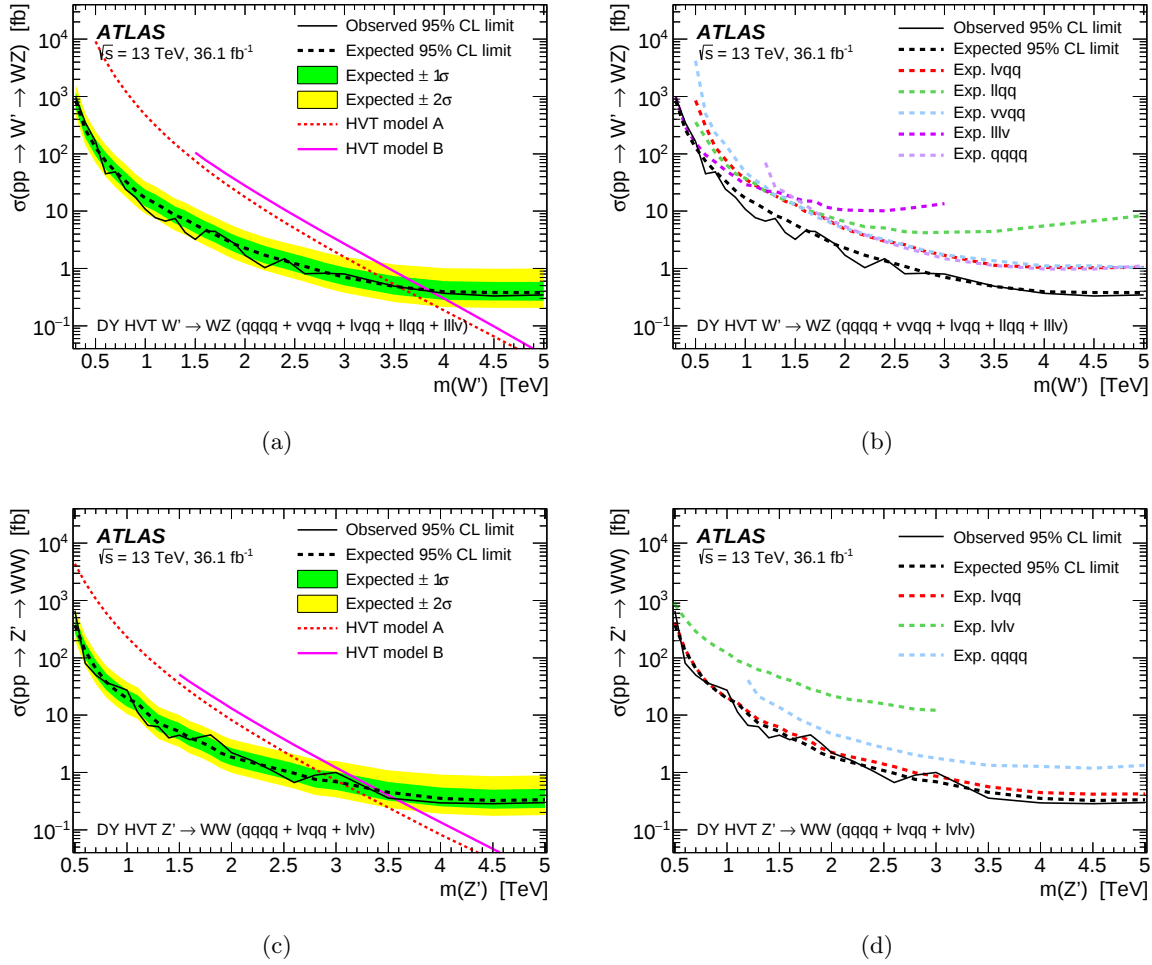


Figure 6.4: The observed and expected 95% CL_s upper limits on the cross-section as a function of mass for the $q\bar{q}$ HVT interpretation in the a) WZ channels and b) WW channels. Overlaid is the theoretically predicted cross-sections for a resonance in this model as a function of mass for two benchmark coupling choices. Figures b) and d) show the expected limits from individual contributing channels.

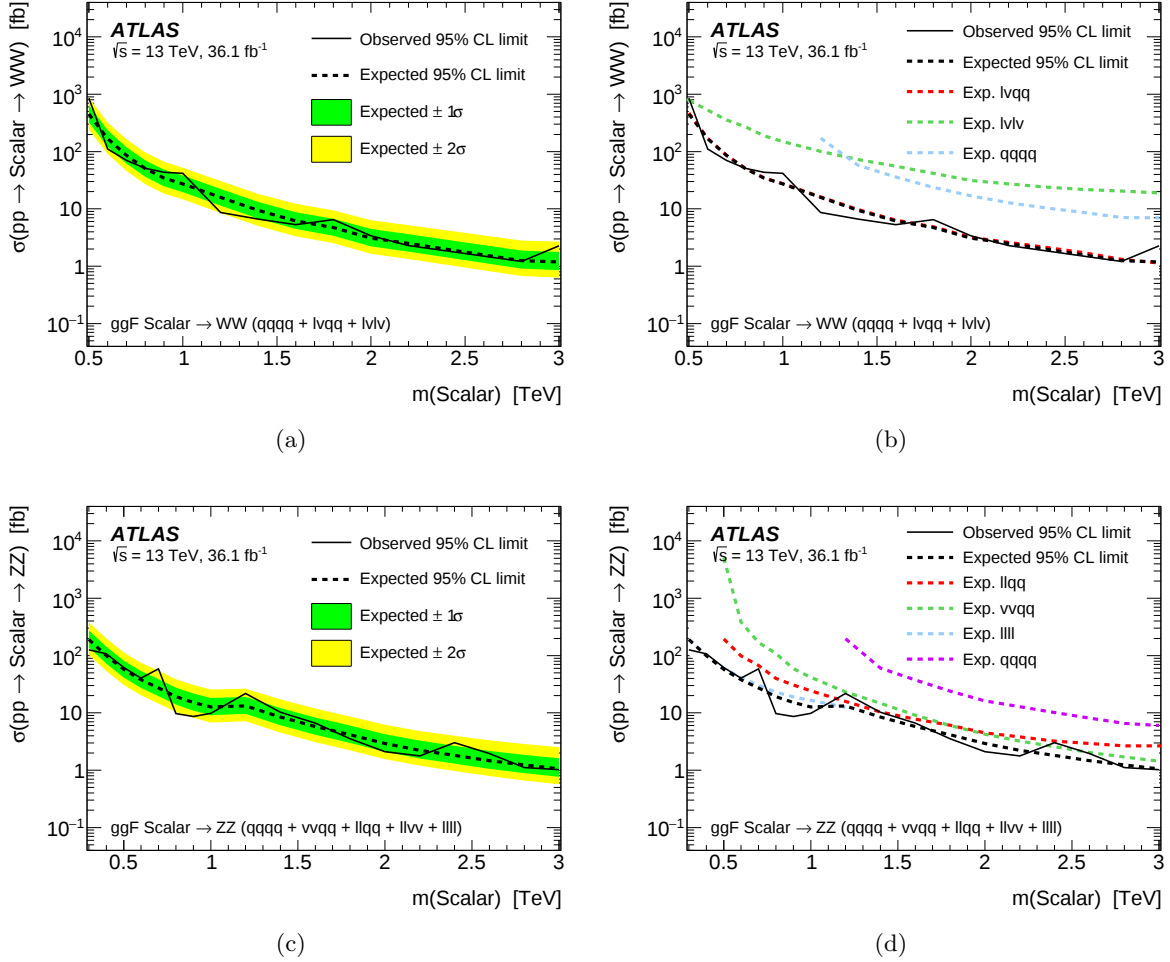


Figure 6.5: The observed and expected 95% CL_s upper limits on the cross-section as a function of mass for the ggF produced heavy Higgs-like scalar interpretation in the a) WW channels and c) ZZ channels. Figures b) and d) show the expected limits from individual contributing channels.

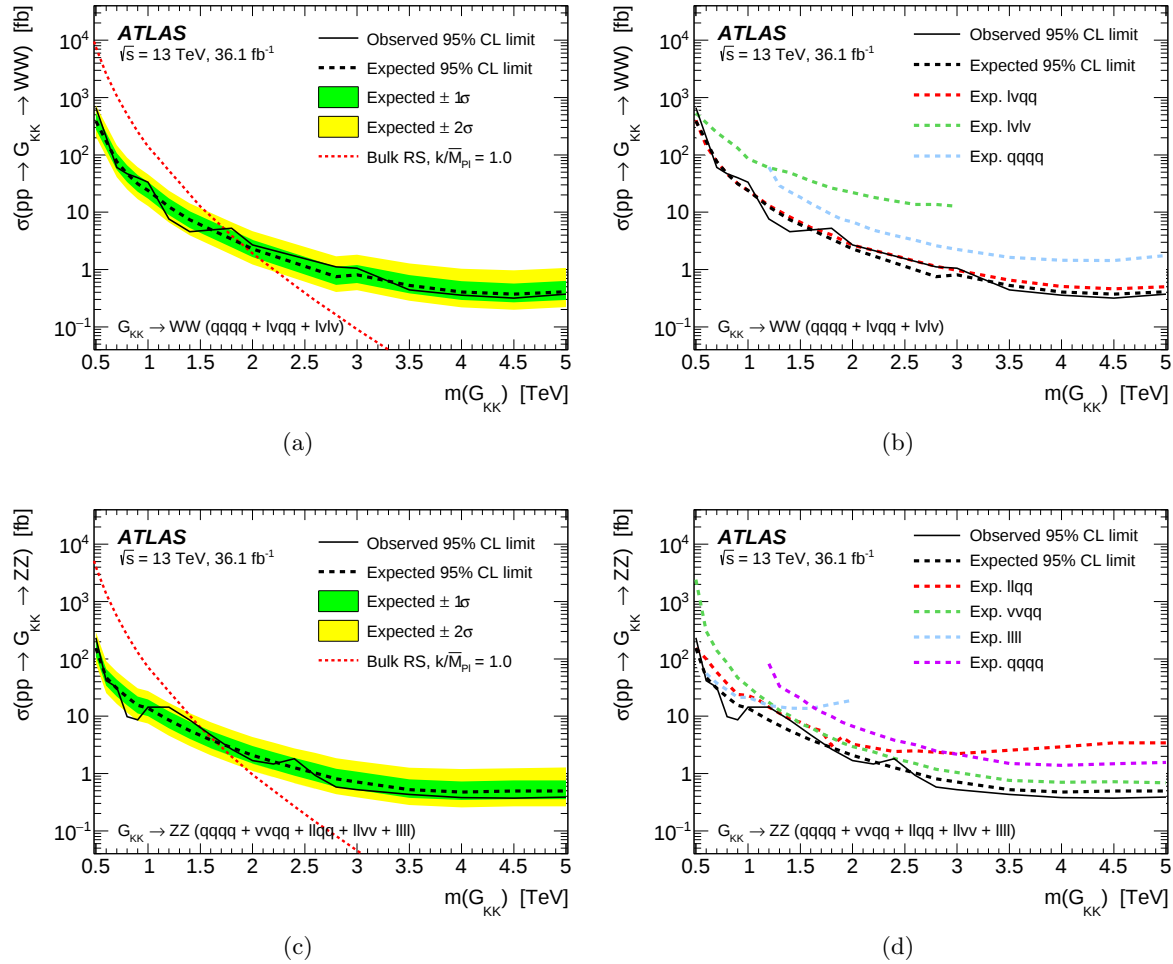


Figure 6.6: The observed and expected 95% CL_s upper limits on the cross-section as a function of mass for the ggF RS graviton interpretation in the a) WW channels and c) ZZ channels. Overlaid is the theoretically predicted cross-sections for a resonance in this model as a function of mass. Figures b) and d) show the expected limits from individual contributing channels.

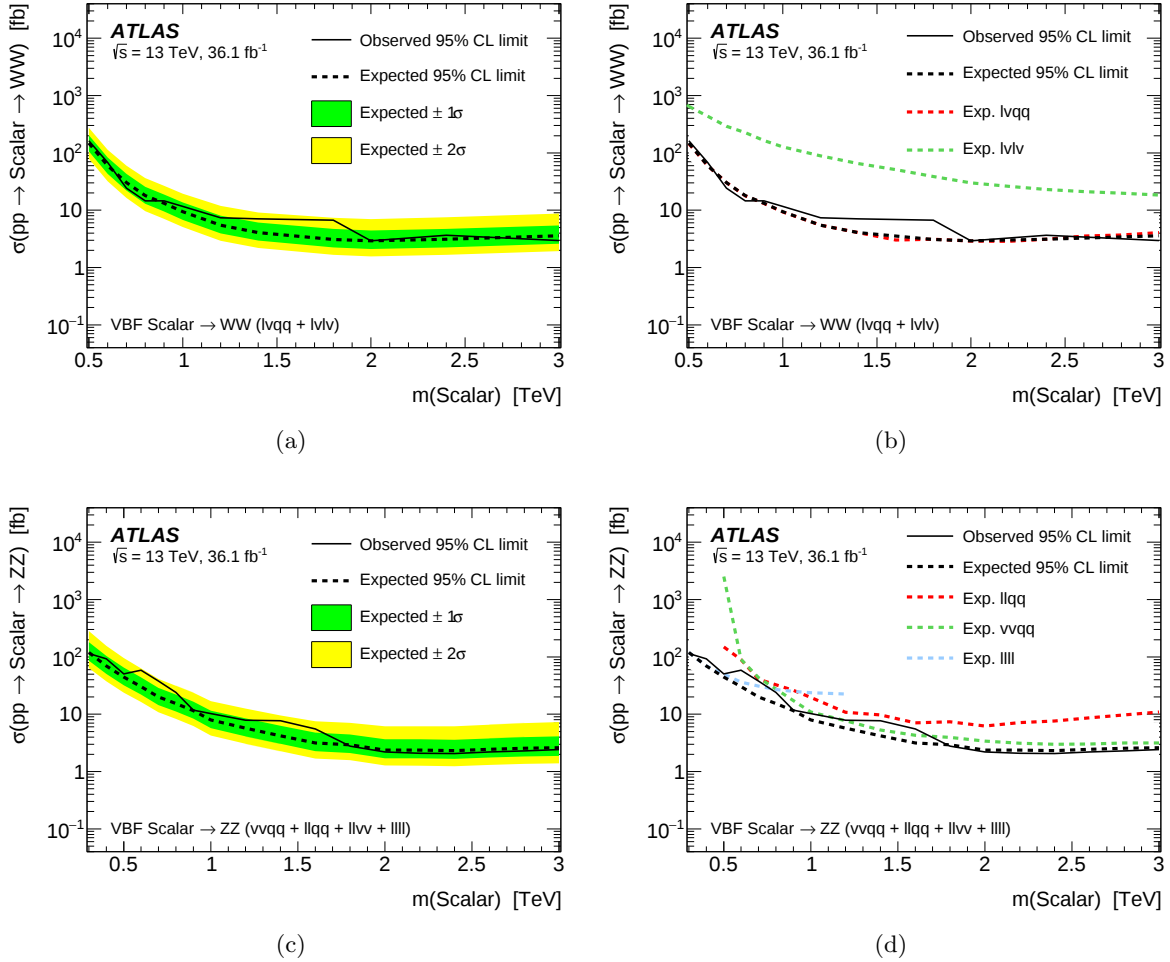


Figure 6.7: The observed and expected 95% CL_s upper limits on the cross-section as a function of mass for the VBF produced heavy Higgs-like scalar interpretation in the a) WW channels and c) ZZ channels. Figures b) and d) show the expected limits from individual contributing channels.

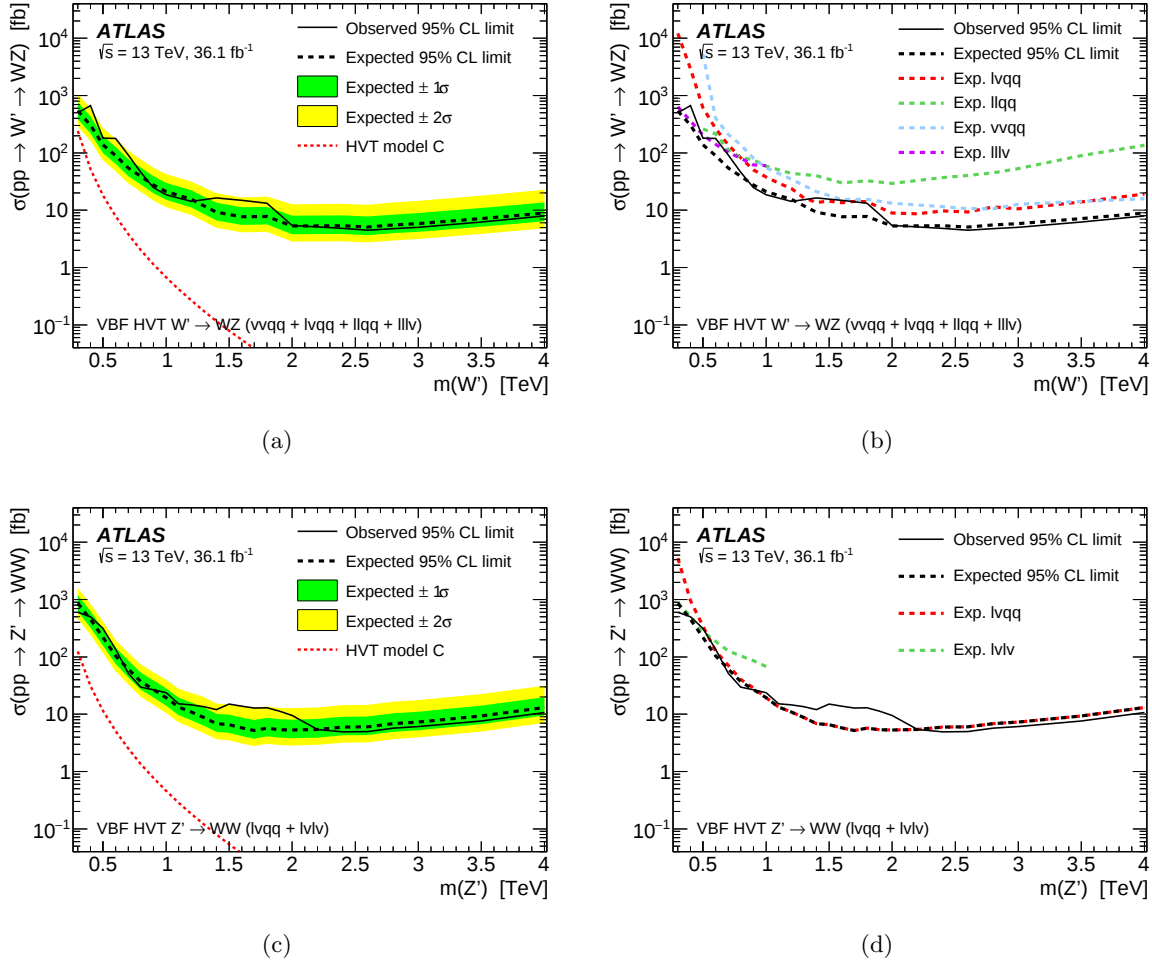


Figure 6.8: The observed and expected 95% CL_s upper limits on the cross-section as a function of mass for the VBF HVT interpretation in the a) WZ channels and c) WW channels. Overlaid is the theoretically predicted cross-sections for a resonance in this model as a function of mass. Figures b) and d) show the expected limits from individual contributing channels.

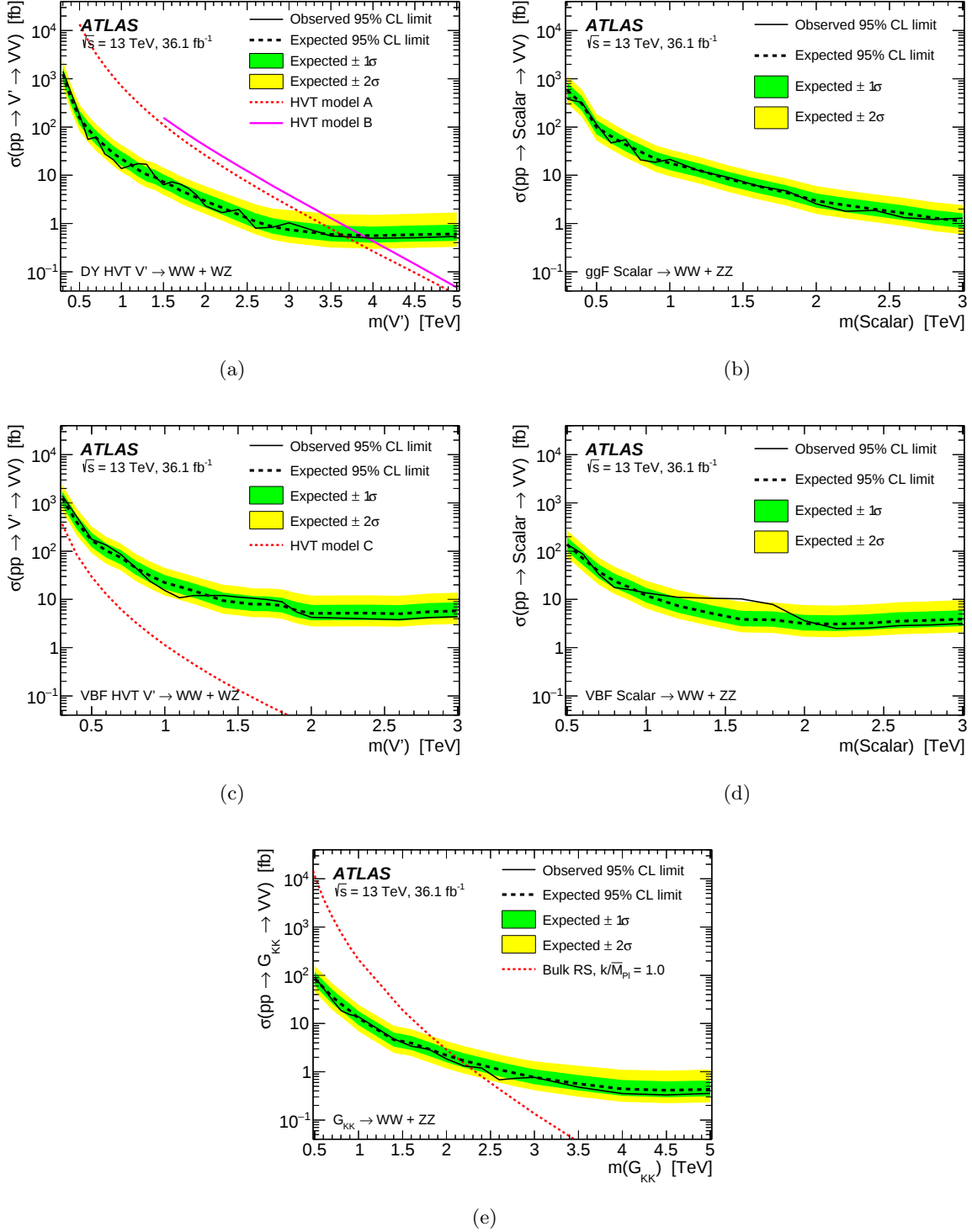


Figure 6.9: The observed and expected 95% CL_s upper limits on the cross-section as a function of mass for the combination of VV channels in the a) $q\bar{q}$ HVT, b) ggF heavy Higgs-like scalar, c) VBF HVT, d) VBF heavy Higgs-like scalar, and e) RS graviton interpretations.

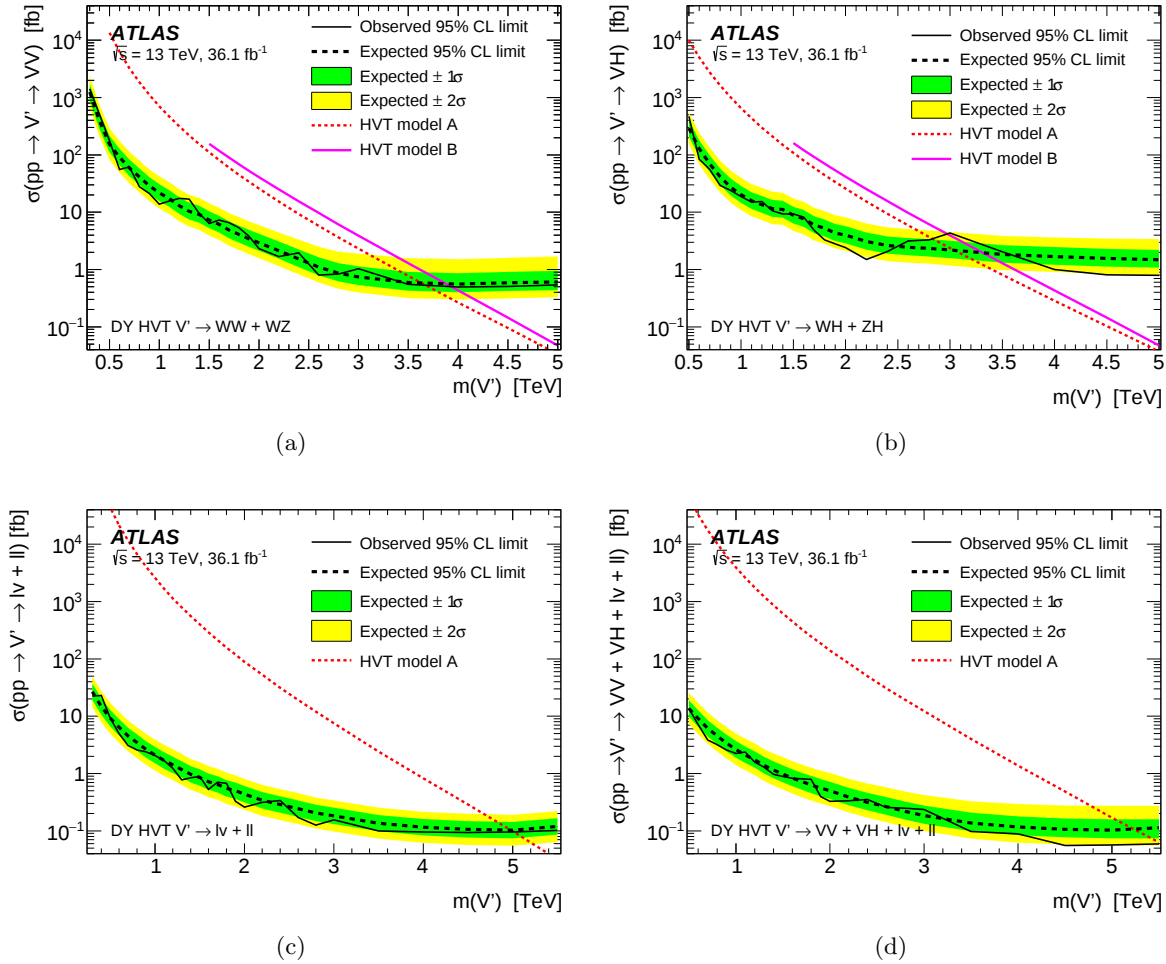


Figure 6.10: The observed and expected 95% CL_s upper limits on the cross-section as a function of mass in the $q\bar{q}$ HVT model for the a) VV combination, b) Vh combination, c) dilepton combination, d) $VV+Vh$ +dilepton combination. The VV combination result from Figure 6.9 is shown again here for comparison.

son+dilepton combination can be seen in Figure 6.10. The Vh channels produce slightly worse limits than the VV limits for both benchmark points, while the dilepton channel has stronger limits in the model A interpretation. The overall diboson and dilepton combination provides a 95% CL_s upper exclusion limit on the HVT W' and Z' mass at approximately 5.5 TeV. While the majority of this sensitivity comes the dilepton results, a 500 GeV improvement on the Model A benchmark mass limit is found by the addition of the diboson channels. It is important to note this is a model dependent statement. The branching ratios of a new resonance to WW , WZ , ZZ , WH , ZH , $\ell\ell$, and $\ell\nu$ are all model dependent. In Figure 6.10d the branching ratios used in the combination are those of the HVT model A predictions. In the HVT model B the coupling to fermions is highly suppressed so the dilepton channels provide effectively no limits. Section 6.5 will describe limits interpretation of the HVT model beyond those of the HVT benchmarks.

6.4.1 Asymptotic vs Toy Limit Comparison

The results shown in Section 6.4 all utilize the asymptotic approximations detailed in Appendix A for evaluation of the limits. These approximations are only valid in the assumption of the large sample limit, which is not typically valid for searches with resonant masses above several TeV. In these ranges the SM background predictions underneath the resonant peak are near zero, which is far from the large sample limit. The asymptotic results are expected to over-estimate the excluded parameter space in such situations.

A proper evaluation of the exclusion limit can be implemented by a Monte-Carlo procedure where the distribution of the test statistic is modeled by rerunning the limit procedure on an ensemble of simulated pseudo-experiments (also known as toy experiments). Since the combination involves many distributions with a high multiplicity of nuisance parameters, the task of running toy simulations for all interpretations and mass points is not computationally feasible.

To measure the expected differences between the asymptotic and full toy procedure in a computationally feasible way, a full toy evaluation was done on a subset of interpretations in a coarse grid of mass points. The ratio of the excluded toy and asymptotic limit was then parameterized with a straight-line as a function of mass. This parameterization was applied as a scale factor on the asymptotic limits.

The result of this procedure can be seen in Figures 6.11 - 6.12 for the combination of HVT $VV+Vh$, and HVT dilepton channels. The asymptotic limits can differ by approximately 50% with respect to the full prediction at 5 TeV mass points. Qualitatively, this indicates that the quoted asymptotic mass limits above are approximately 100 GeV over-optimistic with respect to a full limit calculation. This small difference between the asymptotic and full mass limit calculation can be seen explicitly in Figures 6.11b - 6.12b.

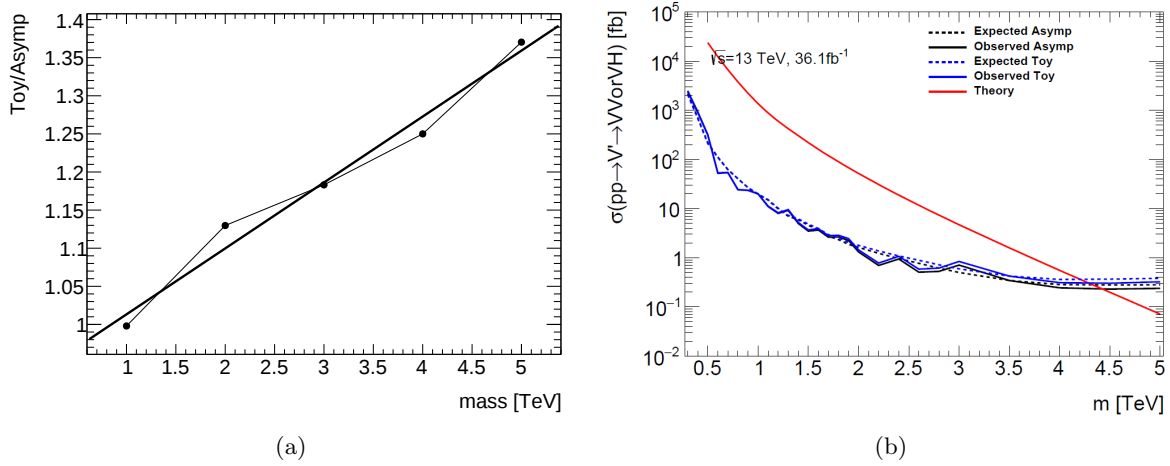


Figure 6.11: a) The ratio of expected limits from full toy Monte-Carlo simulation to asymptotic limits as a function of mass in the HVT $VV+Vh$ fit. b) Comparison of the HVT $VV+Vh$ observed and expected limits (black) from asymptotic approximations and those values scaled to match toy Monte-Carlo predictions (blue). The theoretical prediction for the cross-section is also overlaid (red).

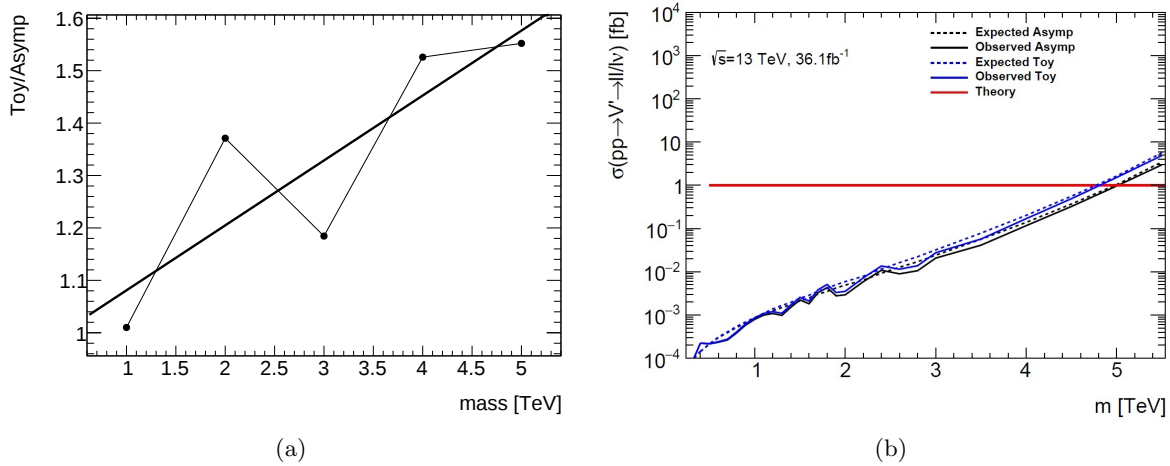


Figure 6.12: a) The ratio of expected limits from full toy Monte-Carlo simulation to asymptotic limits as a function of mass in the HVT dilepton combination. b) Comparison of the HVT dilepton observed and expected limits (black) from asymptotic approximations and those values scaled to match toy Monte-Carlo predictions (blue). The cross-sections are normalized with respect to the HVT model A predictions. The theoretical prediction for the cross-section is also overlaid.

6.5 Exclusion Limits on HVT Coupling Parameters

In Section 6.4, the exclusion limits on the HVT W'/Z' are presented for specific benchmark values of the model parameters. The coupling of the HVT W'/Z' to SM particles is determined solely from independent model parameters in the HVT Lagrangian (Equation 2.30). From just these coupling values and the resonance mass, the cross-section and branching ratios of the new resonance can be analytically calculated. Since all the phenomenology of the HVT model is encompassed by these coupling parameters, we can remove the dependency on benchmark models and instead provide limits directly in the space of coupling parameters.

We will present exclusion limits in two orthogonal coupling planes: the boson g_H v.s. fermion g_f coupling plane (assuming all fermion couplings are equal), as well as the lepton g_l v.s. quark g_q coupling plane assuming $g_H = -0.56$ (the Model A value). For evaluation of these limits, the signal strength is now parameterized by two values, which are the respective coupling being probed. This parameterization allows the contributions of diboson and dilepton channels to vary in the fits across the coupling planes. The branching ratios of W'/Z' to VV/Vh /dilepton are calculated analytically in the coupling planes and the contributing channels scaled by these values in the fit.

The majority of the formalism of Appendix A still applies to multi-dimensional fits. The test statistic is now

$$\lambda(\boldsymbol{\mu}) = \frac{L(\boldsymbol{\mu}|\hat{\boldsymbol{\theta}}(\boldsymbol{\mu}))}{L(\hat{\boldsymbol{\mu}}|\hat{\boldsymbol{\theta}})} \quad (6.1)$$

where $\boldsymbol{\mu}$ is two-dimensional. Similar asymptotic formulae can be derived, but now the test statistic follows a non-central chi-squared distribution with degrees of freedom equal to the dimension of $\boldsymbol{\mu}$. The 95% confidence intervals are calculated in the signal+background hypothesis by the difference of the negative log-likelihood with respect to the minimum of the fit.

Figure 6.13 shows the 95% confidence intervals in the signal+background hypothesis for the g_f and g_H coupling plane from the VV , Vh , and dilepton channels separately. The VV and Vh channels follow a cross pattern as the cross-section to produce the W'/Z' is dependent on the fermion coupling while the branching ratio to the VV/Vh states requires non-zero boson coupling. The dilepton channels on the other hand are almost solely dependent on the fermion coupling which governs both the production and decay rate, with only small sensitivity to the boson coupling due to effects on the total branching ratio.

The combined fit amongst all the channels and a comparison between the individual channels can be seen in Figure 6.14. The diboson channels provide the most stringent constraints at high g_H while providing no limit near $g_H = 0$. The leptonic limits on the other hand have the inverse behavior and offer constraints at $g_H = 0$. The shape of the combined limit is functionally different from the individual diboson and dilepton limits, providing a “kiss-shaped” curve, mostly uniform in g_f . This feature accentuates the power of the combination, as the structure of the limit curves is transformed with respect to the individual diboson or dilepton

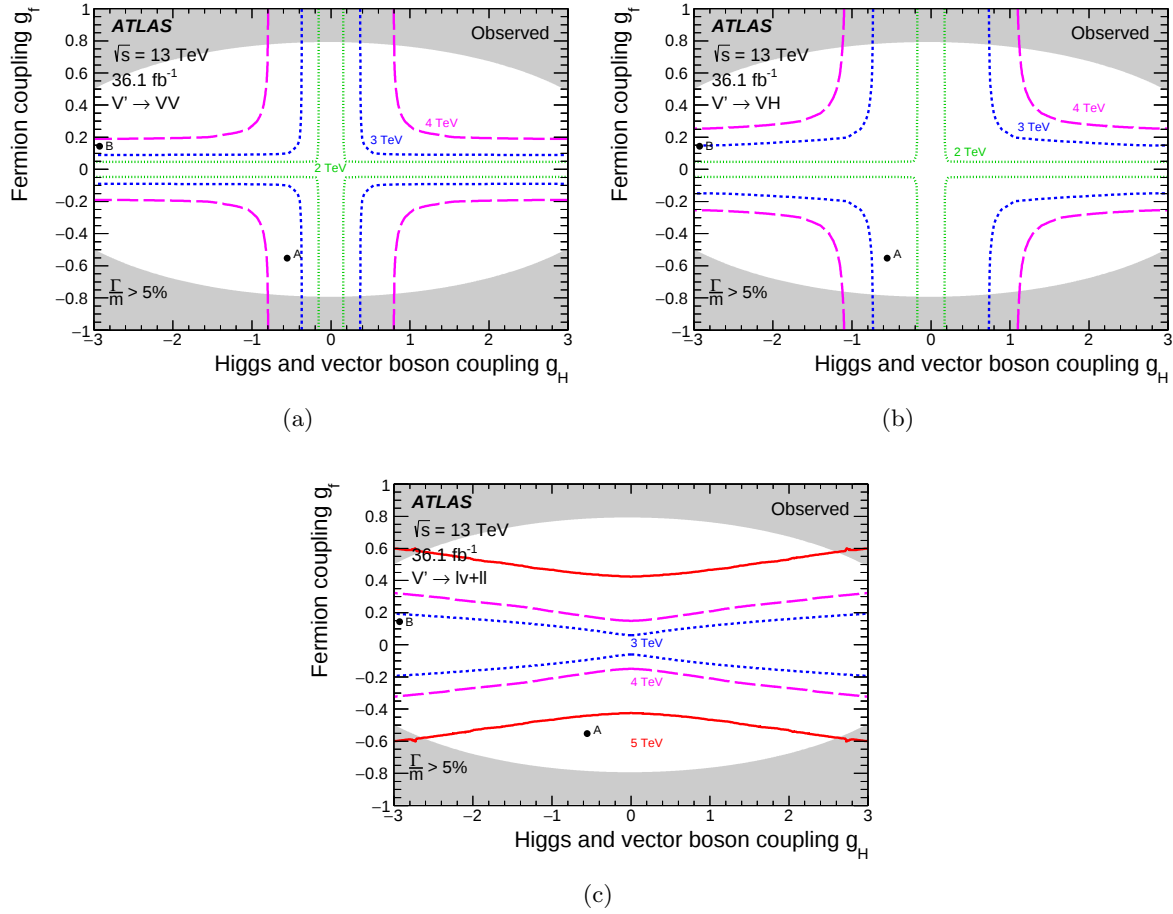


Figure 6.13: Excluded limits at 95% confidence level on the HVT model in the parameter space of fermion coupling and boson coupling for various masses as originating from a) VV , b) Vh , and c) dilepton results. The grey filled area shows the regime where the narrow width approximation used in several of the analysis models is no longer valid. The two points in the plane show the coupling choices of the Model A and Model B benchmark points.

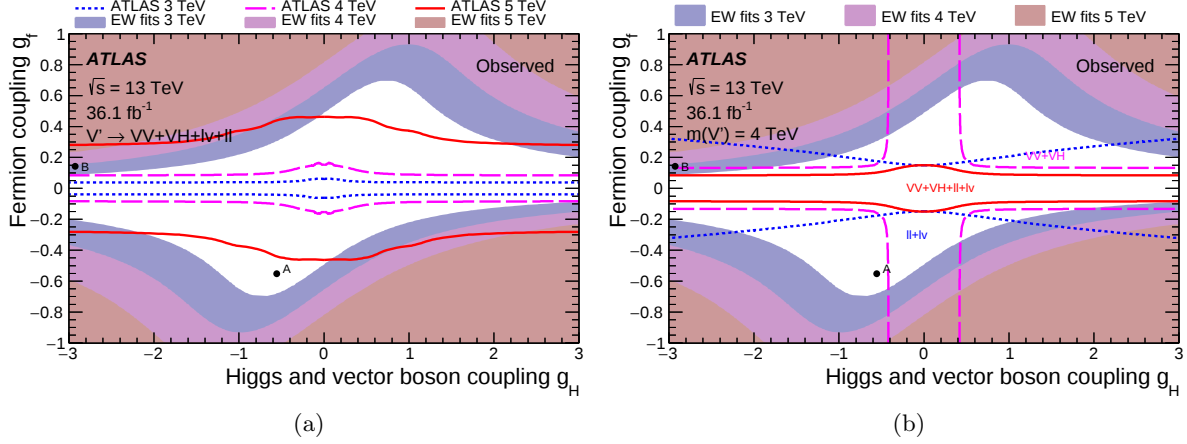


Figure 6.14: Left: The combined diboson and dilepton 95% confidence intervals in the fermion and boson coupling plane for several masses. The corresponding limits from electroweak precision measurements are shown in the filled regions [224]. The two points in the plane show the coupling choices of the Model A and Model B benchmark points. Right: The combined 4 TeV upper exclusion limit (red) and the corresponding limits from the diboson (pink) and dilepton (blue) results only .

analyses. Shown in the filled regions are the limits from electroweak precision measurements described in Ref. [224]. These results include constraints from LEP Z -pole measurements [224], other LEP measurements [225], low-energy experiments [88], top mass measurements [226], and Higgs mass measurements [227]. The combination of ATLAS resonance searches improves on the limits from electroweak limits in almost the whole coupling plane.

The results in the g_l and g_q coupling plane can be seen in Figures 6.15 - 6.16. In this plane the leptonic channels follow a cross-pattern as they are dependent on the quark coupling for production cross-section and leptonic coupling for branching ratio. The diboson channels are mainly only dependent on the quark coupling for production and thus provide limits at $g_l = 0$. This again shows the improvement gained from exploiting complementary searches when casting limits in the parameter plane.

Summary

Searches for new resonantly decaying particles are typically published independently for specific search channels. To extract the most stringent results from a data-set on a model, it is useful to statistically combine such independent searches to one global result. This chapter summarized the combination of all published VV , Vh , and dilepton resonance searches which used 36 fb^{-1} of 13 TeV ATLAS data [9]. These combined results were presented in several steps in order to provide both model agnostic (e.g. $X \rightarrow WW$ combinations) and model dependent results (e.g. $X \rightarrow VV/Vh/\ell\ell$). The HVT model was used to study the diboson and dilepton combinations for spin-1 resonances in a less model-dependent way in which the

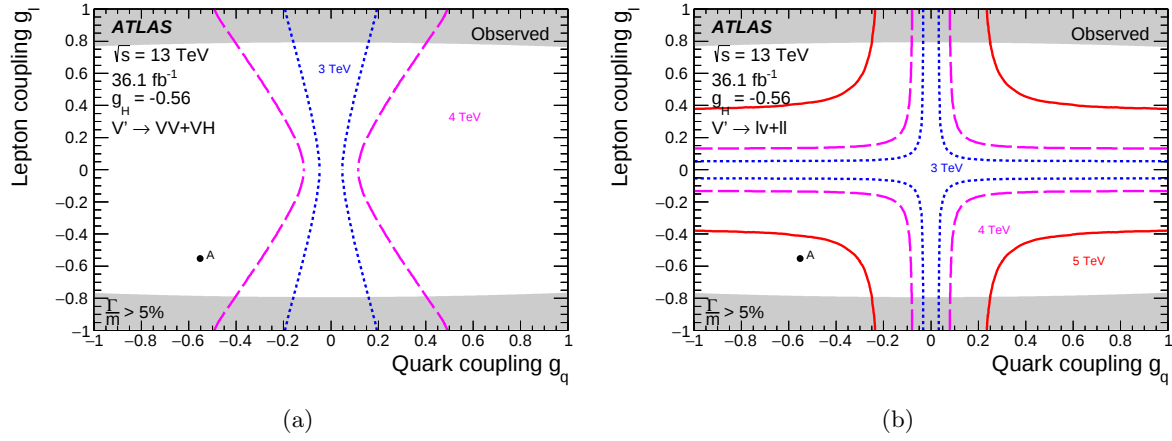


Figure 6.15: Excluded limits at 95% confidence level for the HVT model in the parameter space of lepton coupling and quark coupling for various masses as originating from a) $VV+Vh$, and b) dilepton results. The grey filled area shows the regime where the narrow width approximation used in several of the analysis models is no longer valid. The two points in the plane show the coupling choices of the Model A and Model B benchmark points.

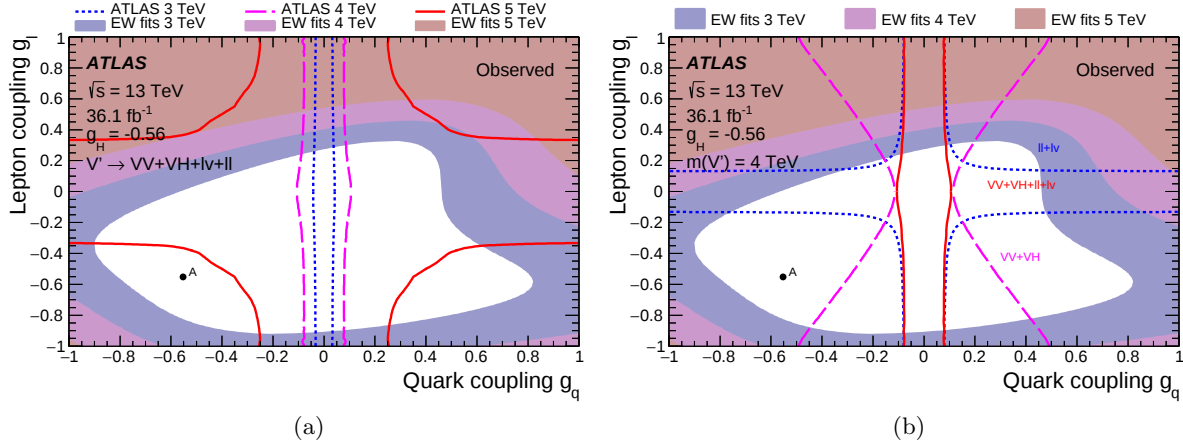


Figure 6.16: Left: The combined diboson and dilepton 95% confidence intervals in the lepton and quark coupling plane for several masses. The corresponding limits from electroweak precision measurements are shown in the filled regions [224]. The two points in the plane show the coupling choices of the Model A and Model B benchmark points. Right: The combined 4 TeV upper exclusion limit (red) and the corresponding limits from the diboson (pink) and dilepton (blue) results only

multi-dimensional coupling plane was directly probed.

The results of this chapter illustrate that the combination of VV channels can improve the 95% CL_s upper excluded limit on the HVT Model A benchmark resonance mass from 2.8 TeV to 3.7 TeV. Further combining with the Vh and dilepton channels increases the limit on the Model A benchmark to 5.5 TeV. When the results are interpreted in the coupling plane of the HVT model, the complementarity of the different channels can be further exploited as different analyses probe different couplings of the HVT model in orthogonal ways. Stringent constraints are then placed on the HVT couplings directly in a form which could not be done by one analysis channel independently.

Chapter 7

Semileptonic VBS search with 36 fb^{-1}

It's a dangerous business, Frodo, going out your door. You step onto the road, and if you don't keep your feet, there's no knowing where you might be swept off to.

— J.R.R. Tolkien, *The Lord of the Rings*

The work presented in Chapters 5 - 6 detailed searches for new physics which decayed resonantly to SM particles. While this is a motivated strategy to search for new physics, it is not the only way. New physics can also appear as gradual deviations in the shape of kinematic observables with respect to the SM predictions. Such signatures could be probed in a model independent way in an EFT framework as discussed in Section 2.3.4. A sketch of how a new physics signal may appear in the context of resonant and non-resonant physics is shown in Figure 7.1. The strategy to search for non-resonant new physics differs from the resonant case as deviations are sought over a wide range of an observable, as opposed to a single localized excess. The search for non-resonant new physics requires an accurate measurement of the SM background.

Chapter 5 optimized an analysis procedure to search for diboson resonance results in the $\ell\nu qq$ channel, with Section 5.2.2 developing additional selection criteria to optimize the search for new particles produced via vector-boson fusion (VBF). The VBF signal interpretation was also implemented in the $\ell\ell qq + \nu\nu qq$ diboson resonance search channels in Ref. [191]. With a reoptimized criteria, the VBF resonance strategy can also be adopted to search and measure the corresponding SM process: electroweak vector-boson scattering (VBS) in the semileptonic channels. In VBS, two gauge bosons radiate from the initial state quarks and scatter off each other. This topology is almost identical to the VBF produced resonances except for the fact the process is not mediated by a new heavy resonance.

VBS is a particularly interesting process for probing the structure and matter content of the SM. As discussed in Section 2.2.4, the VBS process is tightly connected with electroweak symmetry breaking and the Higgs boson. Naively calculating the VBS cross-section without the

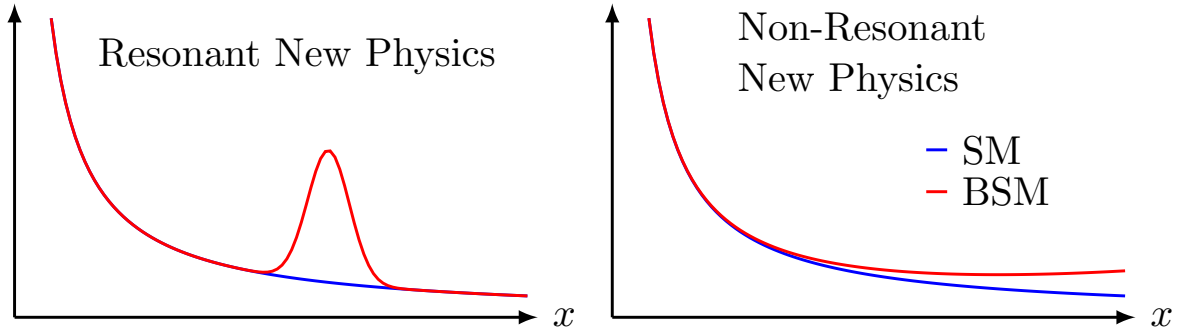


Figure 7.1: Sketch of two possible situation in which new physics can be found. It can appear in the form of new particles which result in increases of the cross-section of an observable at some observable resonant value. New physics can also appear in the form of slight modification of observed kinematics with respect to SM prediction.

Higgs boson, would lead to violation of unitarity at center of mass above approximately 1 TeV. In addition, SM VBS is one of the only processes which involves electroweak quartic interactions at leading order and so can be used to test the gauge structure of the electroweak sector. The other process which electroweak quartic couplings at tree-level is tri-boson production.

SM VBS has already been measured at 6.9 and 5.3σ significance in the leptonic $WW \rightarrow e\nu\mu\nu$ [228] and $WZ \rightarrow \ell\nu\ell\ell$ [229] channels with the ATLAS detector. While the leptonic channels are sensitive to the overall VBS cross-section, searches for anomalous quartic gauge couplings (aQGC) found that the semileptonic channels actually provide more stringent constraint on these new physics operators. This can be seen in Figure 7.2, which shows the upper limits on two aQGC operators from the leptonic and $\ell\nu qq$ VBS searches with 8 TeV ATLAS data [230]. As seen in Table 5.1, the semileptonic channels have larger branching ratios than the leptonic channels which results in increased statistics with respect to those channels. This becomes important in the high energy/momentum tails of the final state object distributions, which is where new physics is still unconstrained. It is then still interesting to measure and observe VBS in the semileptonic channels, even though it is observed in the leptonic channels, as it is more sensitive to changes in observables which can be induced by new physics.

This chapter will summarize the measurement of VBS in the semileptonic channels with 36 fb^{-1} of ATLAS data [10]. This includes all three of the semileptonic channels: $\ell\ell qq$, $\ell\nu qq$, and $\nu\nu qq$. The majority of the analysis procedure is taken directly from the diboson resonance searches in Chapter 5 and Ref. [191] and summarized in Section 7.2. To provide optimal search sensitivity for the VBS process, a Boosted Decision Tree (BDT) algorithm is developed in Section 7.3 to separate signal and background-like events. The BDT distribution will be the final discriminant used to fit the signal and background contribution. Section 7.4 will summarize the relevant systematic uncertainties and in Section 7.5 the results of the search as well as a measurement of the electroweak produced VBS cross-section in the semileptonic channels will be shown.

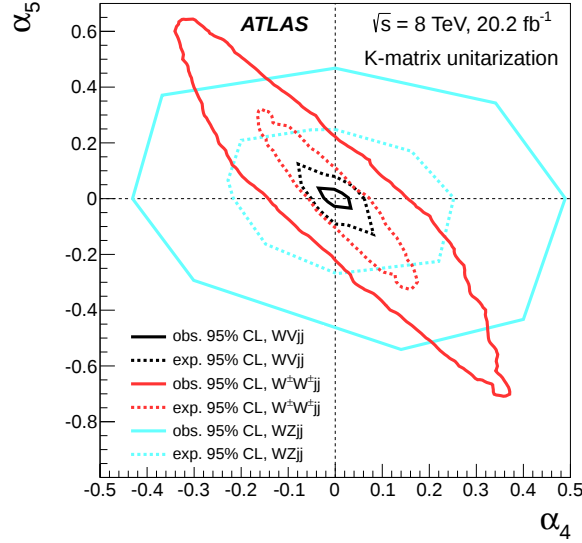


Figure 7.2: Two dimensional 95% confidence intervals on the α_4 and α_5 aQGC operators [231] from the WW leptonic (cyan), WZ leptonic (red), and $WV \ell\nu qq$ channel (black) with 8 TeV ATLAS data. The observed (solid curve) and expected (dashed curve) constraints are shown separately [230].

7.1 Signal and Background Simulation

The MC generation of events with two electroweak gauge bosons and two additional jets, denoted $VVjj$, includes three sets of processes, which are shown in Figure 7.3. The electroweak VBS process is produced at $O(\alpha_{EW}^4)$ and involves triple or quartic gauge couplings. This is the main process of interest for this analysis. At the same order electroweak non-VBS $VVjj$ process can occur, which can not be removed from the VBS diagrams in a gauge invariant way. The electroweak non-VBS contributions will be significantly reduced by the analysis selection, but will still be considered as part of the $VVjj$ signal definition. More than 70% of the events in the WW channel are from $t\bar{t}$ production. The $t\bar{t}$ contribution can be suppressed by 98% by the requirement that no b -tagged jets are within the event. The last set of diagrams are QCD induced processes at $O(\alpha_{EW}^2\alpha_s^2)$, which can be separated at MC generator level and will be similarly reduced during the analysis selection.

The electroweak $VVjj$ processes are generated with MC@NLO v2.4.3 [130] using the NNPD-F30LO [121] set interfaced to Pythia 8.186 [133] for showering and hadronization. The QCD induced $VVjj$ processes are simulated with Sherpa [204]. We will consider the QCD induced contributions as backgrounds and will designate them as SM VV . The MC generators used to model all the relevant backgrounds are the same as those detailed in Section 5.1.2.

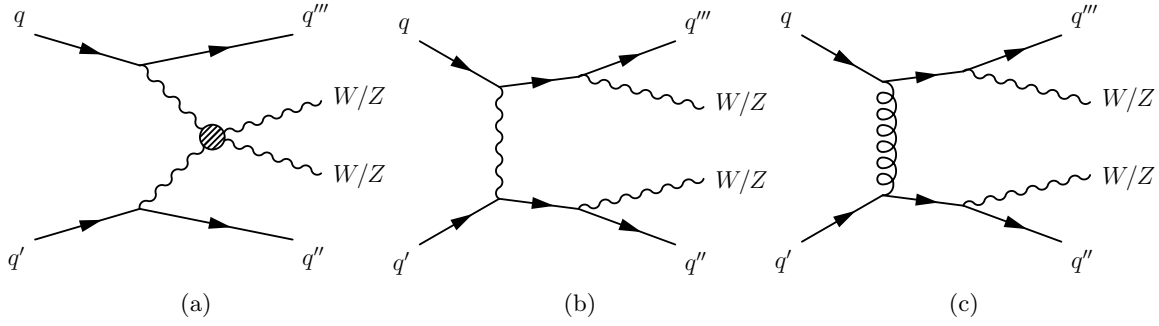


Figure 7.3: Representative Feynman diagrams for the generation of $VVjj$ final state from a) the electroweak VBS processes, b) electroweak non-VBS processes, and c) QCD induced processes.

7.2 Event Selection

The overall analysis strategy is similar to that developed in Section 5.2, but extended to all semileptonic decay modes. The analysis is broken up into three orthogonal channels based on the number of charged leptons¹ in the leptonic W/Z decay: $\ell\ell qq$, $\ell\nu qq$, and $\nu\nu qq$. In each of the three lepton channels a common procedure will be implemented to identify the jets induced by the VBS topology and the hadronic $V_{had} \rightarrow q\bar{q}$ decays. In particular, the $V_{had} \rightarrow q\bar{q}$ decays can be reconstructed as either one large- R jet or two small- R jets, designated the merged and resolved channels. The merged channel is further divided into high-purity (HP) and low-purity (LP) channels based on the output of a hadronic V -tagger. Control regions are defined to measure the V +jets background in each channel, and $t\bar{t}$ background in the $\ell\nu qq$ channel. This scheme results in 9 signal regions and 12 control regions which are fit simultaneously.

With respect to the VBF search of Section 5.2, the VBS-jet and V_{had} jet selection requirements and prioritization scheme were re-optimized simultaneously to provide the best VBS signal sensitivity. All combinations of the following were studied:

- $V_{had} \rightarrow qq$ or VBS jets prioritized
- VBS-jets selected as largest p_T or largest $m(j, j)$ pair
- Resolved V_{had} jets as largest p_T or smallest $|m(j, j) - m_{W/Z}|$ pair
- Boosted V_{had} jet selected as largest p_T , smallest $|m(J) - m_{W/Z}|$, or highest p_T with $|m(J) - m_{W/Z}| < 30$ GeV

The choice which optimized the signal sensitivity was when V_{had} candidate jets were selected first as the pair closest to the W/Z mass, followed by VBS-jets selected as the pair with highest invariant mass. This scheme provides approximately a 12% improvement in the expected sensitivity for the VBS search with respect to the scheme detailed in the resonance search in Section 5.2.2. The requirements on the V_{had} candidates and the VBS-jets are detailed further below. A summary of the event selection is shown in Table 7.1.

¹Only electrons and muons are considered as charged leptons here. The sequential decay of τ -leptons to electrons and muons are still simulated and considered part of the signal and background definition, but direct reconstruction of the τ -lepton is not.

Selection	$\nu\nu qq$	$\ell\nu qq$	$\ell\ell qq$
Trigger	E_T^{miss} triggers	Single-electron triggers Single-muon or E_T^{miss} triggers	Single-lepton triggers
Leptons	0 ‘loose’ leptons with $p_T > 7 \text{ GeV}$	1 ‘tight’ lepton with $p_T > 27 \text{ GeV}$ 0 ‘loose’ leptons with $p_T > 7 \text{ GeV}$	2 ‘loose’ leptons with $p_T > 20 \text{ GeV}$ ≥ 1 lepton with $p_T > 28 \text{ GeV}$
E_T^{miss}	$> 200 \text{ GeV}$	$> 80 \text{ GeV}$	–
$m_{\ell\ell}$	–	–	$83 < m(ee) < 99 \text{ GeV}$ $m(\mu\mu) > -0.0117 \times p_T(\mu\mu) + 85 \text{ GeV}$ $m(\mu\mu) < 0.0185 \times p_T(\mu\mu) + 94 \text{ GeV}$
Small- R jets	$p_T > 20 \text{ GeV}$ if $ \eta < 2.5$, and $p_T > 30 \text{ GeV}$ if $2.5 < \eta < 4.5$		
Large- R jets	$p_T > 200 \text{ GeV}$, $ \eta < 2$		
$V_{had} \rightarrow J$ $V_{had} \rightarrow jj$	V boson tagging, $ m(J) - m_{W/Z} $ $64 < m(jj) < 106 \text{ GeV}$, jj pair with $ m(j_1, j_2) - m_{W/Z} $, leading jet with $p_T > 40 \text{ GeV}$		
VBS-jets	$j \notin V_{had}$, not b -tagged, $\Delta R(J, j) > 1.4$ $\eta(j_1^{VBS}) \cdot \eta(j_2^{VBS}) < 0$, $m(j_1^{VBS}, j_2^{VBS}) > 400 \text{ GeV}$, $p_T > 30 \text{ GeV}$		
Num. of b -jets	–	0	–
Multijet removal	$ \vec{p}_T^{\text{miss}} > 50 \text{ GeV}$ $\Delta\phi(\vec{p}_T^{\text{miss}}, E_T^{\text{miss}}) < \pi/2$ $\min[\Delta\phi(\vec{E}_T^{\text{miss}}, j)] > \pi/6$ $\Delta\phi(\vec{E}_T^{\text{miss}}, V_{had}) > \pi/9$	–	–

Table 7.1: Summary of the event selection in the electroweak VBS search.

7.2.1 Trigger and Lepton Selection

The majority of the event and object selection is the same as in the $\ell\nu qq$ diboson resonance search detailed in Section 5.2. The same list of lowest unrescaled triggers is used to select events in all lepton channels. Single-lepton triggers and E_T^{miss} triggers are used in the $\ell\nu qq$ channels. The $\ell\ell qq$ channel uses the same single-lepton triggers, and the $\nu\nu qq$ channel the same E_T^{miss} triggers as the $\ell\nu qq$ search.

The lepton definitions and selections are the same as those detailed in Section 5.2.1 and in the $\ell\ell qq + \nu\nu qq$ resonance search publication in Ref. [191]. The $\nu\nu qq$ channel requires events with zero leptons passing the loose criteria with $p_T > 7 \text{ GeV}$ and the $\ell\nu qq$ channel requires one tight lepton with $p_T > 27 \text{ GeV}$ and zero additional leptons. In the $\ell\ell qq$ channel, events are required to have two loose leptons of the same flavor with $p_T > 20 \text{ GeV}$, with at least one passing the “medium” identification requirement and $p_T > 28 \text{ GeV}$. Muon pairs are also required to be of opposite charge in the $\ell\ell qq$ channel. The same criterion is not applied in the electron channel due to higher charge misidentification rates. All leptons are required to pass the isolation requirements discussed in Section 5.2.1.

Additional selection requirements on the candidate leptons are defined separately in each lepton channel. In the $\ell\nu qq$ channel, a common requirement of $E_T^{\text{miss}} > 80 \text{ GeV}$ is used to reduce backgrounds. This is a conservative harmonization of the separate resolved and merged channel cuts were used in the resonance search. The full candidate neutrino and $W \rightarrow \ell\nu$ 4-momenta are reconstructed by solving a quadratic equation as detailed in Section 5.2.1. In the $\ell\ell qq$ channel, the invariant mass of the lepton pair is required to be consistent with the Z

mass. The requirement in the electron channel is a flat window of $83 < m(e\bar{e}) < 99$ GeV. In the muon channel the window is $-0.0017 \cdot p_T(\mu\bar{\mu}) + 85 < m(\mu\bar{\mu}) < 0.0185 \cdot p_T(\mu\bar{\mu}) + 94$ GeV in order to recover signal efficiency at high- p_T . This requirement was taken directly from the $\ell\ell qq$ resonance search where it was optimized to give a uniform efficiency across $Z \rightarrow \mu\bar{\mu}$ p_T [191]. In the $\nu\nu qq$ channel a requirement of $E_T^{\text{miss}} > 200$ GeV is required so that events are above the E_T^{miss} trigger turn-on threshold. A significant amount of multi-jet background can still pass the $\nu\nu qq$ selection due to mismeasurements of the jet energies in an event. To reduce this contribution several additional cuts are applied in the $\nu\nu qq$ channel, some of which rely on the track based missing transverse momentum $\vec{p}_T^{\text{miss}}$, which may be accurately measured even if $\vec{E}_T^{\text{miss}}$ is not. The first set of requirements is that the track based missing transverse momentum has $|\vec{p}_T^{\text{miss}}| > 50$ GeV and $\Delta\phi(\vec{p}_T^{\text{miss}}, \vec{E}_T^{\text{miss}}) < \pi/2$. It is also required that the V_{had} candidate and nearest small- R jet have $\Delta\phi(\vec{E}_T^{\text{miss}}, V_{\text{had}}) > \pi/9$ and $\Delta\phi(\vec{E}_T^{\text{miss}}, j) > \pi/6$ respectively. The multi-jet contribution is negligible with respect to the other backgrounds after application of these cuts in the $\nu\nu qq$ channel.

7.2.2 Signal Jet Selection

The reconstruction of the hadronically decaying boson follows the same procedure as was done in the resonant search: either as a single large- R jet, known as the merged channel, or as two small- R jets, known as the resolved channel. Large- R jets are required to pass the same selection criteria and V -tagger detailed in Section 5.2.3. The merged channels are separated into merged high-purity (HP) and merged low-purity (LP) regions if they pass the 50% V -tagger efficiency working point, or fail the 50% working point but pass the 80% working point. Events are not separated into hadronic W or Z channels. If a large- R jet passes both the W -tagger and Z -tagger requirements, the one with the smallest $|m(J) - m_{W/Z}|$ is used. The collection of track-jets (detailed in Section 4.3.2) are used to identify possible b -hadrons within the large- R jet. Track-jets are b -tagged and ghost associated to large- R jets in the event. The number of associated track-jets and whether they are b -tagged are used as an input to the BDT algorithm in Section 7.3 for some of the lepton channels.

The resolved channel follows a slightly different procedure than the diboson resonance search in Section 5.2.4. The $V_{\text{had}} \rightarrow q\bar{q}$ small- R jet pair is selected as the two jets with smallest $|m(j_1, j_2) - m_{W/Z}|$. This pairing strategy is chosen as it provides the best signal efficiency and is discussed further in context of the VBS-tag jets in Section 7.2.3. To suppress backgrounds after the dijet pair selection, the leading jet is required to have $p_T > 40$ GeV and the dijet mass within $64 < |m(j_1, j_2) - m_{W/Z}| < 106$ GeV. This window is a merger of the W and Z windows used in Section 5.2.4.

For both the resolved and merged $\ell\nu qq$ channel, events are removed if a single b -tagged jet is found. As was done in the resonance search, if an event passes the merged selection it is removed from consideration of the resolved selection.

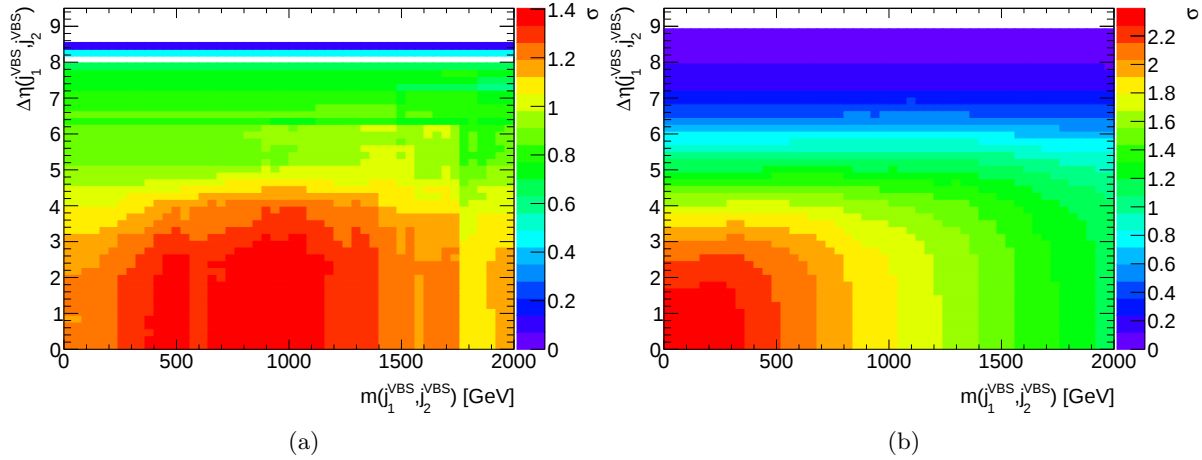


Figure 7.4: The expected binned significance of the VBS signal in the $\ell\nu qq$ a) merged and b) resolved channel as of cuts on the invariant mass and $\Delta\eta$ during jet selection.

7.2.3 VBS Jet Selection

After the selection of the V_{had} candidate, the event is searched for two energetic and forward jets which are indicative of the VBS topology (j_1^{VBS} and j_2^{VBS}). If an event is in the merged selection, only small- R jets with $\Delta R > 1.4$ from the large- R jet are considered possible VBS-tag jets. If an event is in the resolved selection, the two jets which make the $V_{had} \rightarrow q\bar{q}$ candidate are removed from consideration. Similar to Section 5.2.2, all jet pairings which are in opposite hemispheres and are not b -tagged are considered. If multiple jets satisfy this criteria, the pair with highest $m(j_1^{VBS}, j_2^{VBS})$ is chosen.

The invariant mass and opening angle requirements on the VBS-jet pair were also re-optimized with respect to those of the resonance search. The signal sensitivity as defined in Equation 5.1 can be seen in Figure 7.4 as a function of these cuts in the merged and resolved $\ell\nu qq$ channels. It can be seen that a wide-range of cut thresholds almost equally optimize the signal sensitivity. In particular, it is found that no $\Delta\eta(j_1^{VBS}, j_2^{VBS})$ is explicitly needed and $m(j_1^{VBS}, j_2^{VBS}) > 400 \text{ GeV}$ is a compromise between optimal resolved and merged channels values. It is also found that applying a requirement on the VBS-jet p_T of 30 GeV after the VBS-jet selection improves the background rejection as opposed to applying it during the jet selection as in the resonant search Section 5.2.2.

7.2.4 Background Determination

Sideband regions are defined using the same strategy as Sections 5.2.5 - 5.2.6 to constrain the major backgrounds during the final fit. In the $\ell\nu qq$ channel, a W +jets enriched control region (WCR) and top control region (TCR) are defined identically as in Section 5.2.5 as the events which pass all the selections except for inverted cuts on the V_{had} mass and b -jet veto respectively. In the $\ell\ell qq$ channel, a Z +jets control region (ZCR) is defined with the same inverted V_{had} cut

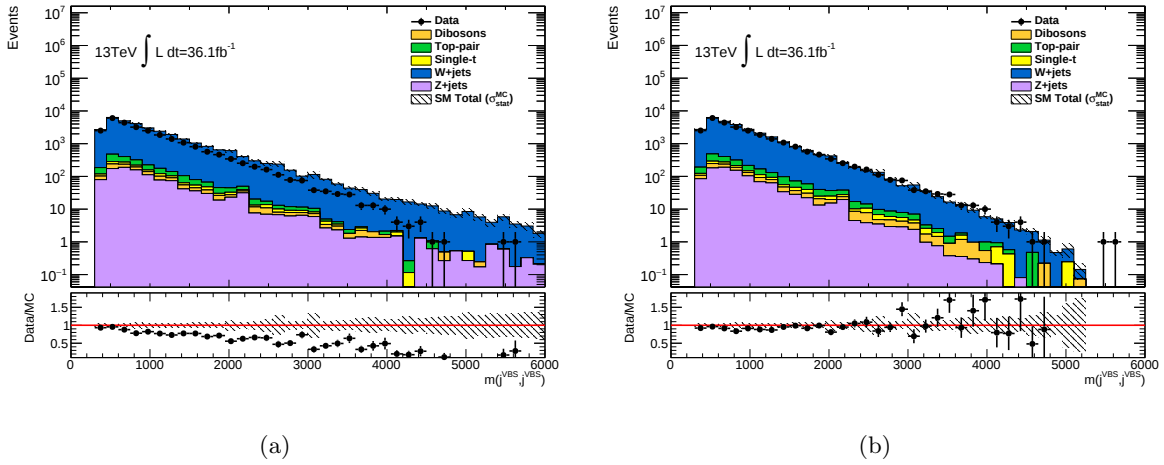


Figure 7.5: The predicted $m(j_1^{VBS}, j_2^{VBS})$ MC distributions in the resolved WCR channel a) before and b) after the reweighting procedure. The bottom subplot shows the data/MC ratio in each bin. More than a 50% difference in MC prediction and data can be found above 3 TeV. After the reweighting procedure the data/MC ratio is found to be uniform. The background normalization is extracted during the fit.

as the WCR. The $\nu\nu qq$ channel has significant contributions from both Z +jets and W +jets, so the corresponding mass sideband region is designated a V +jets control region (VCR). Separate control regions are defined for the merged HP, merged LP, and resolved channels. The ZCR is more than 95% pure in Z +jets in all regions, while the WCR is 86% and 77% pure in the merged and resolved regions. The respective TCR is 79% and 59% pure in $t\bar{t}$ backgrounds.

A mismodelling of the Sherpa V +jets MC samples in the $m(j_1^{VBS}, j_2^{VBS})$ distributions is observed in this analysis. This can be seen in the resolved WCR region shown in Figure 7.5 where MC can predict more than twice the yields of data above $m(j_1^{VBS}, j_2^{VBS}) > 3$ TeV. To correct for this effect, the V +jets MC is re-weighted in a data driven way. In the $\ell\nu qq$ WCR and $\ell\ell qq$ ZCR the data/MC ratio of the $m(j_1^{VBS}, j_2^{VBS})$ variable is linearly parameterized after removing the non- V +jets contributions. The resulting linear function is used as a re-weighting factor for the W +jets and Z +jets in the respective signal regions. These are also applied in the $\nu\nu qq$ signal regions. To account for selection biases when extrapolating these factors to the $\nu\nu qq$ channel, the difference in the shape of W +jets and Z +jets MC in $\nu\nu qq$ with respect to $\ell\nu qq$ and $\ell\ell qq$ is also parameterized and applied as an additional re-weighting factor. After the re-weighting, good data/MC agreement is found in all signal and control regions. The re-weighting procedure is found to have no noticeable effect on observables uncorrelated with $m(j_1^{VBS}, j_2^{VBS})$. We will consider a conservative estimate of the V +jets modelling uncertainty in Section 7.4.

The shape of all backgrounds is taken directly from MC and the normalization of the $t\bar{t}$ and V +jets are extracted from the control regions during the fit. A 1% multi-jet background contribution is expected in the resolved $\ell\nu qq$ channel and is estimated with the same “fake-

factor” method described in Section 5.2.5.

7.3 BDT Optimization

To optimize the search sensitivity and exploit correlations between different observables, a classifier to distinguish the VBS signal from the background is designed using a boosted decision trees (BDT) algorithm. The BDT response will be used as the final observable in the fit, and will not be directly cut on. The training and validation is done using the BDT module of the TMVA package [155]. The BDT was separately optimized for both merged and resolved channels in all three lepton channels. The inputs provided to the BDT algorithm were the VBS signal MC and the sum of background MC in the signal region. For conciseness, only the optimization in the $\ell\nu qq$ channel will be discussed below as the $\ell\ell qq$ and $\nu\nu qq$ channels followed the same procedure.

The first step of the optimization identified the variables which provided the largest separation power between signal and background. The separation of an observable x is defined as [155]

$$S = \frac{1}{2} \int \frac{(f_s(x) - f_b(x))^2}{f_s(x) + f_b(x)} dx \quad (7.1)$$

where $f_s(x)$ and $f_b(x)$ are the signal and background probability distributions, which are estimated by binned distributions. Over 50 input variables were tested, including all the base (e.g. ℓ) and compound object (e.g. $W \rightarrow \ell\nu$) 4-momenta, and the relative $\Delta\eta/\Delta\phi$ angles between them. Other observables considered include the number of tracks with $p_T > 500 \text{ MeV}$ associated to the jet (n_{trk}), and the jet width defined as

$$w(j) = \frac{\sum_k \Delta R(j, j_k) p_T(j_k)}{p_T(j)} \quad (7.2)$$

where k represents an individual EMTopo cluster associated to the jet with 4-momenta j_k . The boson centrality defined as

$$\zeta_V = \min(\Delta\eta_+, \Delta\eta_-) \quad (7.3)$$

where

$$\begin{aligned} \Delta\eta_+ &= \max(\eta(j_1^{VBS}), \eta(j_2^{VBS})) - \max(\eta(W), \eta(V_{had})) \\ \Delta\eta_- &= \min(\eta(W), \eta(V_{had})) - \min(\eta(j_1^{VBS}), \eta(j_2^{VBS})) \end{aligned} \quad (7.4)$$

was also tested.

Only the top subset of weakly-correlated inputs with the highest separation were considered in the BDT training. One-by-one the variables with highest separation were added into the BDT training until no further improvement on the BDT classification was found or over-training was observed. For evaluation of the BDT performance, half the MC samples were used in the

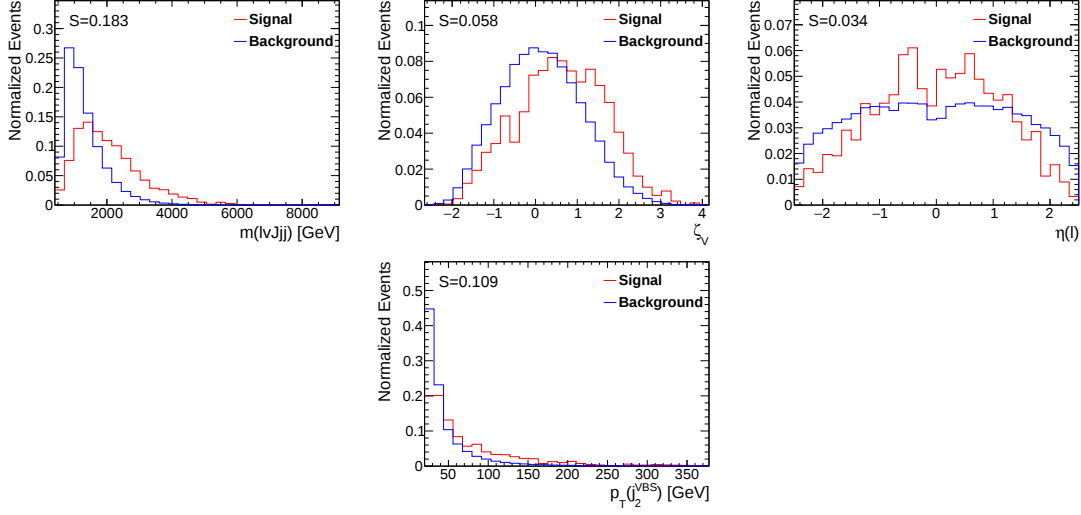


Figure 7.6: The signal (red) and background (blue) distributions of the input variables used for the merged $\ell\nu qq$ BDT. The separation S is shown for each variable.

training and the other half for validation. The distributions of the final input variables for the $\ell\nu qq$ BDT are shown in Figures 7.6 - 7.7 for the merged and resolved channels. The distributions of all input variables are found to be well modeled by MC predictions.

Due to limited signal statistics in the VBS MC samples, the training was done in an extended signal region where the $m(j_1^{VBS}, j_2^{VBS}) > 400$ GeV cut is removed. In the merged channels the requirement on the V -tagging is also removed. In order to not bias the BDT, variables strongly correlated to the VBS-jet mass and large- R jet mass/ $D_2^{\beta=1}$ are removed from the training. The BDT trained in the extended region is found to have similar discrimination power when applied in the full signal region definitions.

The hyper-parameters and architecture of the BDT algorithm are also optimized. It was found for this analysis that using a Gradient-Boosted BDT algorithm outperforms the conventional AdaBoost algorithm [155]. This is expected since Gradient-Boosting is known to perform better in situations with low sample statistics. After use of Gradient-Boosting, the BDT performance was found to be invariant for reasonable choices of hyper-parameters.

Additional checks on the stability of the BDT training were also performed. The split of samples into training/validation is an arbitrary decision and so two choices were tested: splitting on odd/even event numbers, or with respect to the first half and second half of the sample. Both splittings resulted in BDT distributions which were consistent within statistical uncertainty. To provide further stability and increase the effective statistics used, a k -fold cross-validation technique is implemented where the output BDT distribution is averaged across BDT scores trained on sub-samples with training and validation events switched. The BDT distributions of one of the k -fold trained samples is shown in Figure 7.8 for the merged and resolved $\ell\nu qq$ channel. This illustrates the separation power provided by the BDT in these regions.

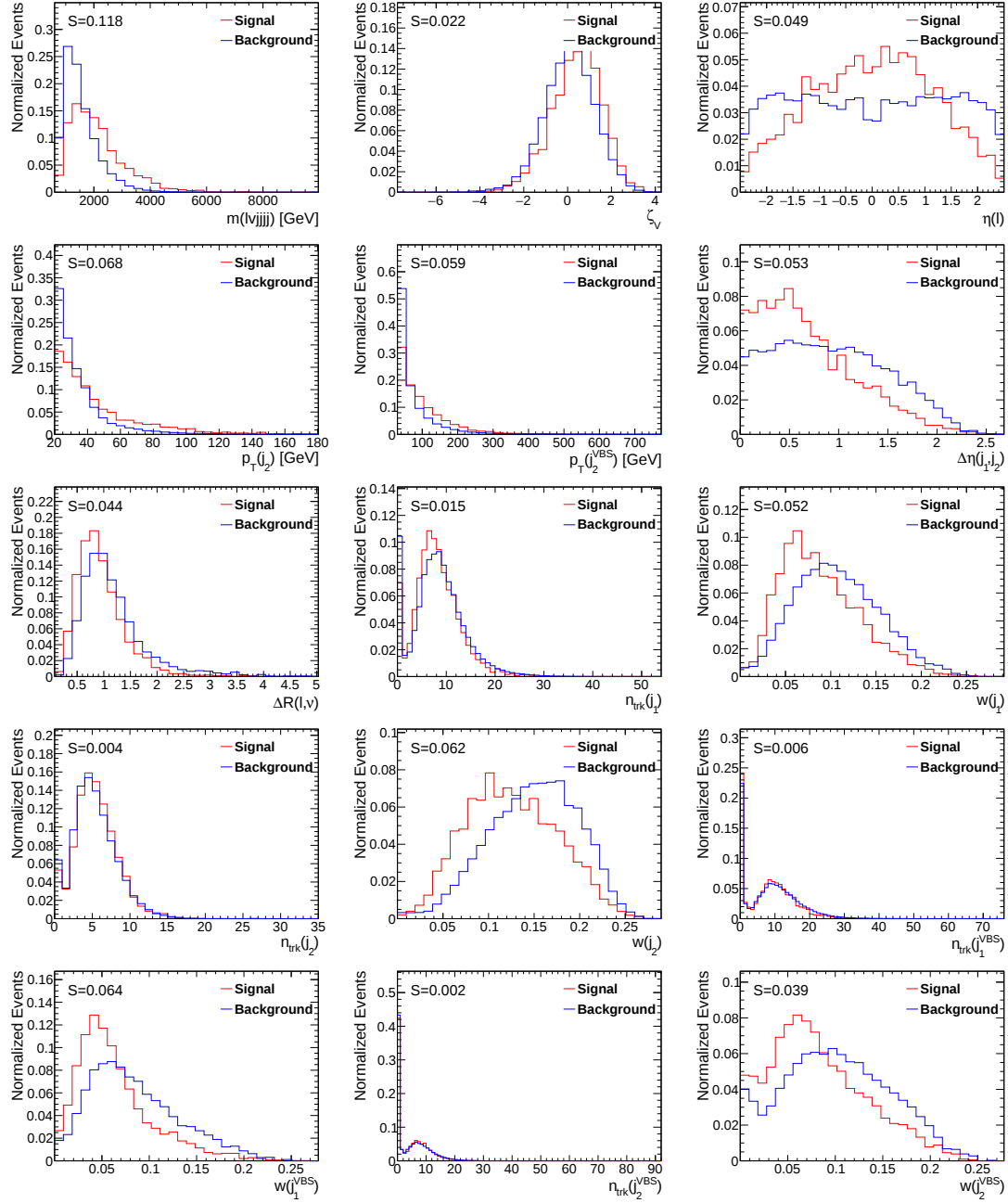


Figure 7.7: The input signal (red) and background (blue) distributions of the input variables used for the resolved $\ell\nu qq$ BDT. The separation S is shown for each variable.

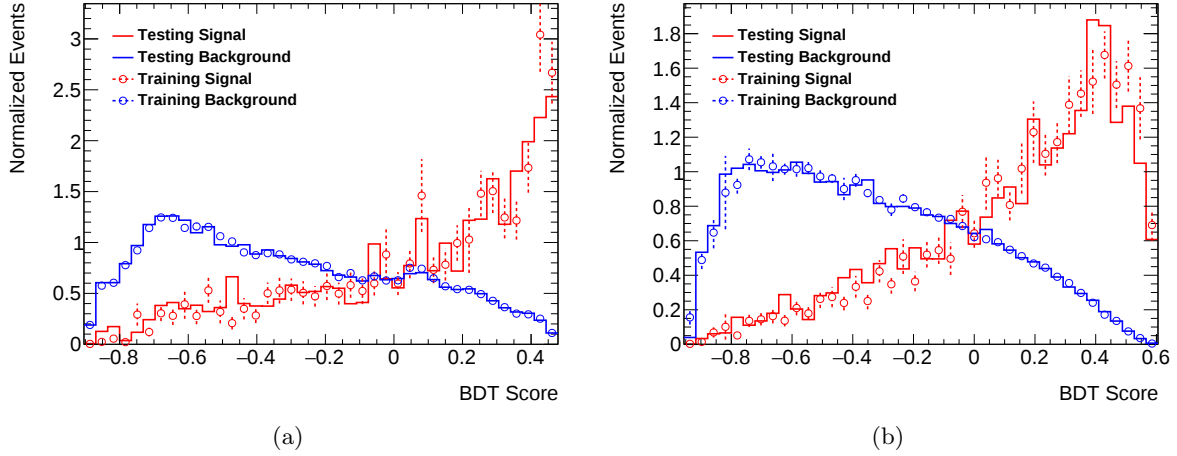


Figure 7.8: The $\ell\nu qq$ a) merged and b) resolved BDT score distributions for the signal (red) and background (blue) as evaluated during the final stage of training (points) and on an uncorrelated validation sample (curves). Statistical uncertainty is shown for the training sample distribution.

7.4 Systematics

All sources of uncertainty considered in the diboson resonance search, detailed in Section 5.3, except for the resonance signal uncertainties, are common to this analysis as well. For conciseness, only new sources of uncertainty will be discussed in this section. As a reminder, normalization uncertainties are quoted as percentages differences of the affected sample only, and not with respect to the overall background.

The only new experimental sources of uncertainty considered in this analysis are those from the data-driven V +jets Sherpa correction detailed in Section 7.2.4. The $m(j_1^{VBS}, j_2^{VBS})$ variable is used as one of the inputs in the $\ell\ell qq$ and $\nu\nu qq$ channel BDT, so the impact of this uncertainty on the fit results can be large. To estimate the uncertainty in a data-driven way, an initial conservative estimation is provided by the two-point variation between the re-weighted and original distributions. Since an overly conservative estimate is provided, during the profile-likelihood fit (see Section A for details) the size of this uncertainty is allowed to converge to the value which provides the best fit to the data. The data then provides an indirect measurement of the uncertainty from the constraint of this nuisance parameter in the fit. An additional uncertainty on the extrapolation of the ZCR and WCR V +jets measurements into the $\nu\nu qq$ channel is estimated by the double ratio of Sherpa and Madgraph V +jets yields between the $\nu\nu qq$ channel and the other two lepton channels. This uncertainty is applied as an additional Gaussian normalization factor in the $\nu\nu qq$ signal regions with variance of 8 (14)% for W +jets events and 22 (42)% for Z +jets events in the merged (resolved) signal region.

In addition to the experimental systematics, several new modeling systematics are considered. Following the same procedure as for the V +jets modeling uncertainties evaluated in

Section 5.3.2, the normalization uncertainty on the QCD-induced SM VV samples is estimated from variations in the PDF set and scale choices. The normalization uncertainty is estimated to be approximately 20%. This value is inflated to 30% to be conservative as the Sherpa matrix element calculation is NLO only at 1 and 2 additional partons for this process.

The uncertainty on the electroweak $VVjj$ signal normalization from missing higher order QCD terms is estimated by varying renormalization, factorization, and α_s scales in the calculation. The PDF uncertainty is estimated by the envelope of the uncertainties of the nominal NNPDF [121] and the MMHT [123] PDF set. The uncertainty from ISR/FSR modeling is estimated by comparison with alternative Pythia [200] tunes. The contributions of each of these uncertainties varies per signal region, but are approximately 2% for the QCD scale, 5% for PDF, and 4% from ISR/FSR uncertainties. The effect of the missing QCD-electroweak interference terms in the MC generation is measured by simulating all contributions with MC@NLO v2.4.3. The effect on the overall normalization is found to be under a percent, with shape differences on the truth $m(j_1^{VBS}, j_2^{VBS})$ distributions below 6% (10%) in the resolved (merged) regions.

7.5 Results

This analysis uses the statistical framework described in Appendix A to model the data with a binned profile-likelihood. In each of the fit regions, binned distributions for the signal and background are evaluated from the MC predictions and fit to the observed data. The normalization of the electroweak $VVjj$ process is parameterized by the ratio of observed to expected cross-section $\mu = \frac{\sigma^{obs}}{\sigma^{exp}}$, which is left to float to the best value during the fit. The sources of uncertainty are treated as additional nuisance parameters in the fit which are contained by external measurements. The W +jets, Z +jets, and $t\bar{t}$ normalizations are unconstrained parameters, and allowed to converge to the values which provide optimal agreement with the data. All systematics and normalizations are correlated amongst all regions, except for the V +jets and $t\bar{t}$ normalizations which are separated between the merged and resolved regions.

All 21 regions described in Section 7.2.4 are fit simultaneously. This includes 9 signal regions for the $\ell\ell qq$, $\ell\nu qq$, and $\nu\nu qq$ channels in each of the merged HP, merged LP and resolved regions. To provide optimal signal sensitivity, the specifically trained BDT response for that region is the fit observable in the signal regions. The remaining regions are the WCR, ZCR, VCR, TCR control regions for each of the merged HP, merged LP, and resolved channels. For the WCR, ZCR, and VCR regions, $m(j_1^{VBS}, j_2^{VBS})$ is used in the fit in order to constrain the V +jets reweighting uncertainty. These regions are also used to extract the V +jets normalization. The TCR is only used to extract the $t\bar{t}$ normalization.

The results are presented as a search for the electroweak $VVjj$ production in Section 7.5.1 and as a cross-section measurement in a fiducial region in Section 7.5.2.

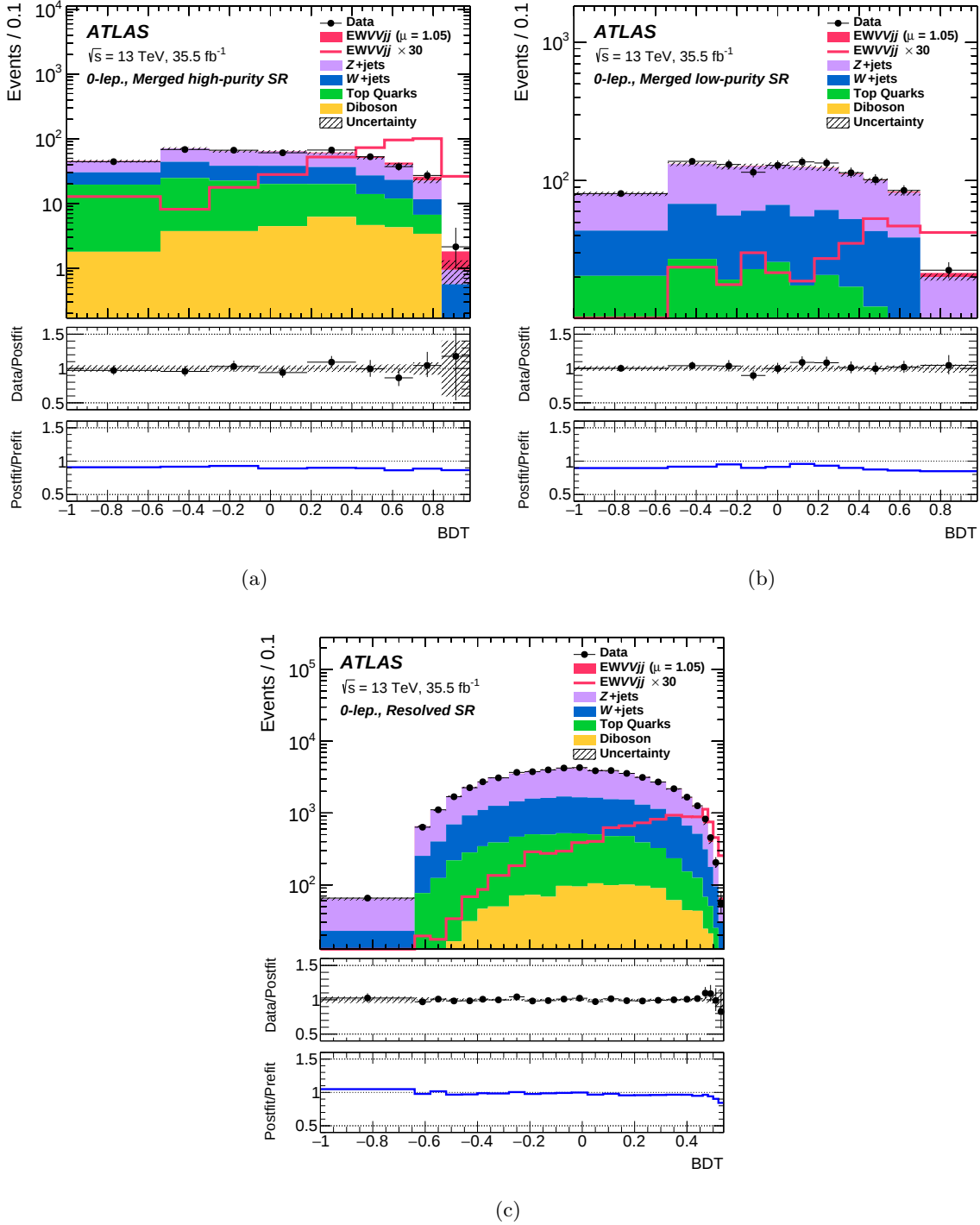


Figure 7.9: The BDT distributions in the $\nu\nu qq$ a) merged HP, b) merged LP, and c) resolved signal regions after the simultaneous final fit to all regions. The data is shown in black and the filled regions show the separated contributions of the measured background processes. The total uncertainty, including both systematics and statistical uncertainties is shown as black hashed region on the background envelope. The filled red histograms show the contribution of the measured electroweak $VVjj$ process stacked over the background MC contribution, and the overlaid red curve shows just this contribution scaled by a factor of 30 for improved visibility. The middle sub-plots show the ratio of the data to MC in each bin of the fit. The bottom sub-plots show the ratio in each bin of the post-fit MC to pre-fit MC yield

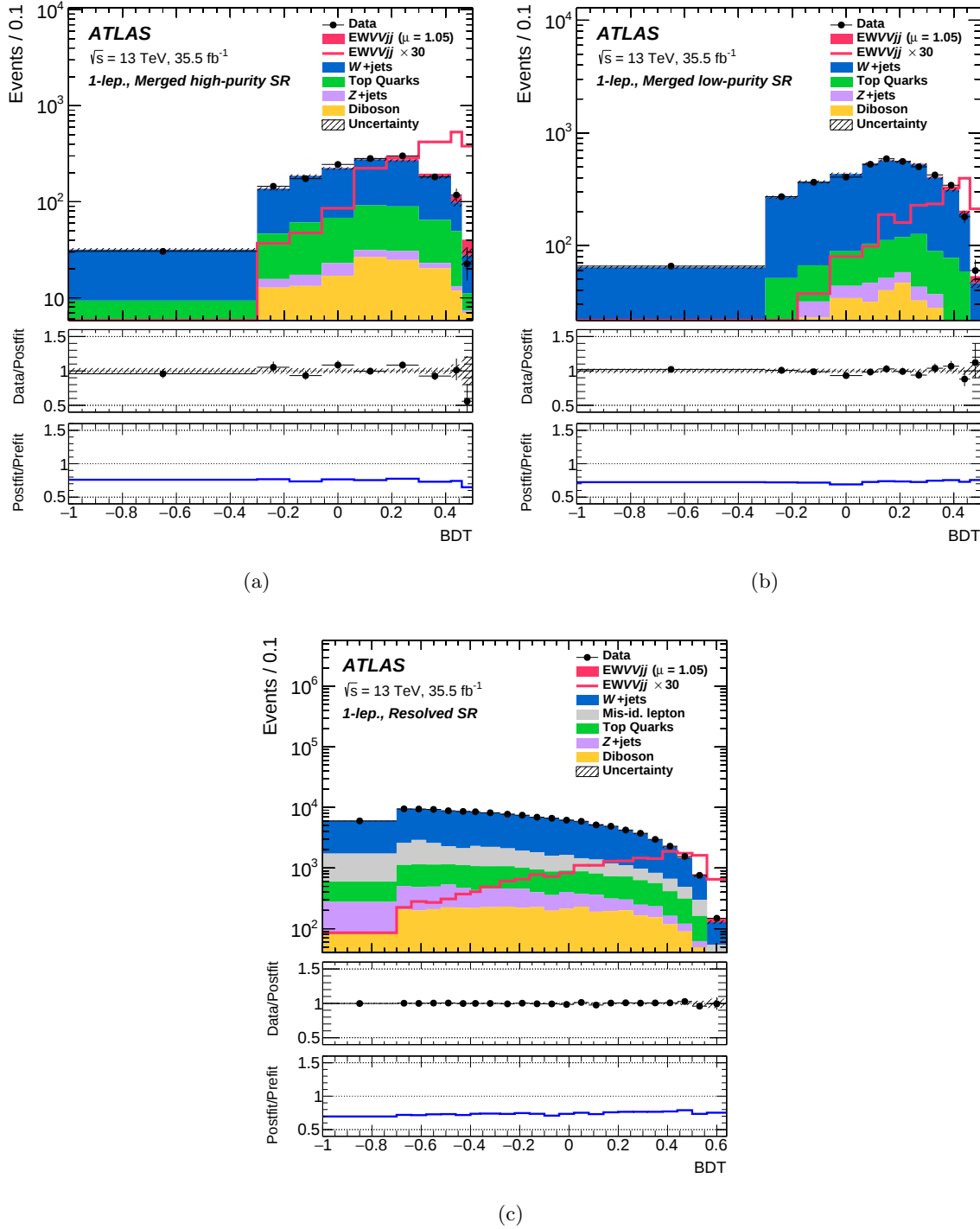


Figure 7.10: The BDT distributions in the $\ell\nu qq$ a) merged HP, b) merged LP, and c) resolved signal regions after the simultaneous final fit to all regions. The data is shown in black and the filled regions show the separated contributions of the measured background processes. The total uncertainty, including both systematics and statistical uncertainties is shown as black hashed region on the background envelope. The filled red histograms show the contribution of the measured electroweak VVjj process stacked over the background MC contribution, and the overlaid red curve shows just this contribution scaled by a factor of 30 for improved visibility. The middle sub-plots show the ratio of the data to MC in each bin of the fit. The bottom sub-plots show the ratio in each bin of the post-fit MC to pre-fit MC yield

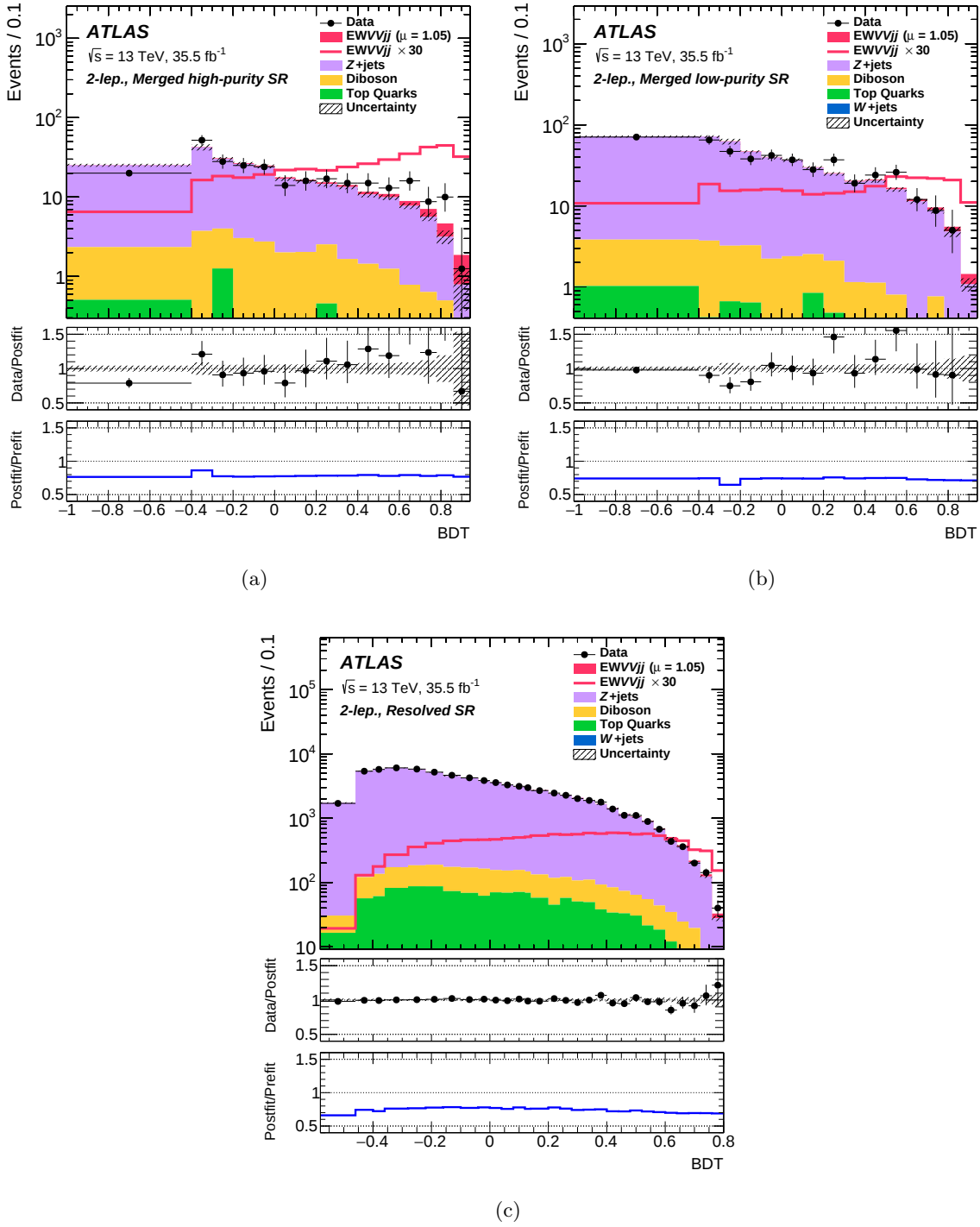


Figure 7.11: The BDT distributions in the $\ell\ell qq$ a) merged HP, b) merged LP, and c) resolved signal regions after the simultaneous final fit to all regions. The data is shown in black and the filled regions show the separated contributions of the measured background processes. The total uncertainty, including both systematics and statistical uncertainties is shown as black hashed region on the background envelope. The filled red histograms show the contribution of the measured electroweak VVjj process stacked over the background MC contribution, and the overlaid red curve shows just this contribution scaled by a factor of 30 for improved visibility. The middle sub-plots show the ratio of the data to MC in each bin of the fit. The bottom sub-plots show the ratio in each bin of the post-fit MC to pre-fit MC yield

7.5.1 Search for Electroweak $VVjj$ Production

The distributions of the BDT score in the signal regions after the simultaneous fit to all regions can be seen in Figures 7.9 - 7.11. The data and fit agree after the inclusion of the electroweak $VVjj$ component. The measured background only p -value corresponds to a significance of 2.7σ , in comparison to the 2.5σ expected. The value of 2.7σ is slightly below the conventional limit of 3σ to claim “evidence” for the signal plus background hypothesis². We therefore make no claim that the data excludes the background only hypothesis.

If the BDT is replaced with the best single variable discriminant, the diboson invariant mass $m(VV)$, the expected background only p -value is 1.45σ . This highlights the importance of the development of the BDT discriminant in this analysis, which increased the search sensitivity by a relative 70%.

The best fit result of the measured electroweak $VVjj$ signal strength with respect to the SM prediction is

$$\mu = 1.05^{+0.42}_{-0.40} = 1.05 \pm 0.20(stat)^{+0.37}_{-0.34}(syst) \quad (7.5)$$

with the statistical and systematic uncertainties displayed separately and as combined values. The data is most consistent with the signal+background hypothesis, with signal strength near the SM prediction. The shape of the negative log-likelihood of the fit around the minimum, and a breakdown of the fitted signal strength in individual lepton channels is displayed in Figure 7.12.

The relative contributions of the top ranked sources of uncertainty on the final signal strength value are summarized in Table 7.2. Both statistical and systematic uncertainties are relevant contributions to the final measurement. The largest systematics uncertainties are the limited generated MC statistics and the MC modeling uncertainties. Future analyses can improve on this result by utilizing both a larger data set and by increasing the relative size of the generated MC samples.

7.5.2 Fiducial Cross-section Measurement

Even though we do not claim “evidence” for the $VVjj$ signal, we find the data is more consistent with the signal+background hypothesis than the background only hypothesis. We therefore still use the data to extract a measurement of the $VVjj$ cross-sections, which can be useful for comparison with other experimental results and theoretical predictions.

The cross-section measurement is presented in the region of phase space set by the detector and kinematic acceptance of this analysis, denoted as the fiducial region. This region is defined at the level of “truth” MC predictions (defined in Section 4.1) that match the full reconstruction level analysis as closely as possible. A fiducial measurement is made as opposed to an inclusive

²Technically p -values are only claims on the null (background only) hypothesis, and make no claim that the signal plus background hypothesis is correct. We will maintain the definition of 3σ as “evidence” and 5σ as “discovery” of the corresponding SM signal process to follow the typical convention of the literature.

Uncertainty source	σ_μ
Total uncertainty	0.41
Statistical	0.20
Systematic	0.35
Theoretical and modeling uncertainties	
MC statistics	0.17
Z +jets	0.13
Floating normalizations	0.09
W +jets	0.09
Diboson	0.09
Signal	0.07
$t\bar{t}$	0.06
Multijet	0.04
Experimental uncertainties	
Large- R jets	0.08
b -tagging	0.07
Small- R jets	0.06
E_T^{miss}	0.04
Pileup	0.04
Luminosity	0.03
Leptons	0.02

Table 7.2: The symmetrized uncertainty σ_μ on the best-fit signal-strength parameter for the top ranking sources of uncertainties. The top ranking modeling and experimental systematics are shown separately. The “Floating normalizations” term includes the uncertainty of the V +jets, and $t\bar{t}$ normalizations as extracted from data in the control region.

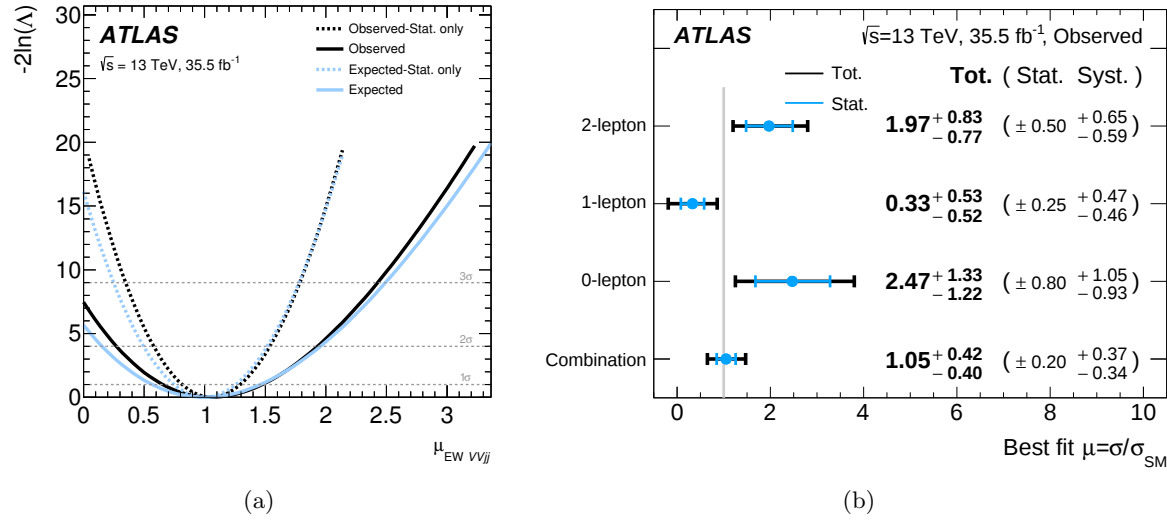


Figure 7.12: a) A scan of two times the negative log-likelihood of the fit around the minimum as a function of the signal strength μ . The observed (black) and expected (blue) scan can be seen including all systematics (solid) and solely the statistical systematics (dashed). Gray dashed lines indicate threshold values for evaluating confidence intervals on the signal strength. b) The fitted value of the signal strength and confidence intervals as evaluated from just statistical uncertainties (blue) or with all systematics (black) from fits to individual semileptonic channels and the combination of them.

Object selection		
Leptons	$p_{\text{T}} > 7 \text{ GeV}, \eta < 2.5$	
Small- R jets	$p_{\text{T}} > 20 \text{ GeV}$ if $ \eta < 2.5$, and $p_{\text{T}} > 30 \text{ GeV}$ if $2.5 < \eta < 4.5$	
Large- R jets	$p_{\text{T}} > 200 \text{ GeV}, \eta < 2.0$	
Event selection		
Leptonic V selection	$\nu\nu qq$	Zero leptons, $p_{\text{T}}(\nu\nu) > 200 \text{ GeV}$
	$\ell\nu qq$	One lepton with $p_{\text{T}} > 27 \text{ GeV}$, $p_{\text{T}}(\ell\nu) > 80 \text{ GeV}$
	$\ell\ell qq$	Two leptons, with leading (subleading) lepton $p_{\text{T}} > 28$ (20) GeV
		$83 < m_{\ell\ell} < 99 \text{ GeV}$
Hadronic V selection	Merged	One large- R jet, $\min(m(J) - m_{W/Z})$ $64 < m(J) < 106 \text{ GeV}$
	Resolved	Two small- R jets, $\min(m(j_1, j_2) - m_{W/Z})$ $p_{\text{T}}(j_1) > 40 \text{ GeV}, p_{\text{T}}(j_2) > 20 \text{ GeV}$ $64 < m(j_1, j_2) < 106 \text{ GeV}$
	Tagging-jets	Two small- R non- b jets, $\eta(j_1^{VBS}) \cdot \eta(j_2^{VBS}) < 0$, highest $m(j_1^{VBS}, j_2^{VBS})$ $m(j_1^{VBS}, j_2^{VBS}) > 400 \text{ GeV}, p_{\text{T}}(j_{1,2}^{VBS}) > 30 \text{ GeV}$
Number of b -jets	$\nu\nu qq$	–
	$\ell\nu qq$	0
	$\ell\ell qq$	–

Table 7.3: Definition of the fiducial phase-space used for the measurement of electroweak $VVjj$ cross-section.

Fiducial phase space	Expected $\sigma^{fid,exp}(VVjj)$ [fb]	Measured $\sigma^{fid,obs}(VVjj)$ [fb]
Merged	11.4 ± 0.7 (theo.)	12.7 ± 3.8 (stat.) $^{+4.8}_{-4.2}$ (syst.)
Resolved	31.6 ± 1.8 (theo.)	26.5 ± 8.2 (stat.) $^{+17.4}_{-17.1}$ (syst.)
Inclusive	43.0 ± 2.4 (theo.)	45.1 ± 8.6 (stat.) $^{+15.9}_{-14.6}$ (syst.)

Table 7.4: Expected and measured fiducial level cross-sections when allowing the merged and resolved channels to be fit with separate signal strengths or a single common one.

measurement to facilitate direct comparison of different MC generator predictions. The fiducial regions are defined in Table 7.3, where merged HP and LP are combined into one region.

Since the fiducial region is defined at the MC generator level, it is not directly comparable to the data since it does not include detector and reconstruction effects. To account for this, the difference between the fiducial and reconstruction level analysis is quantified using a single “correction factor” defined as

$$C = \frac{N(\text{reco} : \text{analysis requirements})}{N(\text{truth} : \text{fiducial requirements})} \quad (7.6)$$

where N is the number of events in the category as estimated by MC. Here “truth” denotes the MC generator’s prediction at the generator output level, while “reco” denotes the results after application of detector simulation and object reconstruction. C is then an overall normalization factor to be applied in the fiducial region which accounts for detector efficiency, acceptance, and reconstruction performance. The fiducial level cross-section is then be measured as

$$\sigma^{fid,obs}(VVjj) = \mu' \sigma^{fid,exp}(VVjj) = \mu' \frac{N_{exp}}{L \cdot C} \quad (7.7)$$

where N_{exp} is the expected truth-level MC yield of the $VVjj$ sample in the fiducial region, and L is the integrated luminosity. To provide a cross-section measurement for comparison with other generators, it is useful to separate the signal modeling uncertainties from the measurement result. The signal modeling uncertainties are therefore removed from the fit, and a new μ' is extracted from the data. The signal modeling uncertainties are instead presented as theoretical uncertainties on the expected measurement. The expected cross-sections are those predicted with MC@NLO 2.4.3 at leading order.

Two sets of fiducial measurements are presented. In Table 7.4 the measured and expected fiducial cross-sections are shown when allowing merged and resolved regions to be fit with separate normalization factors or when they are fit with a common signal strength. All lepton channels use common signal strengths in these fits. The results of the fiducial cross-section measurement when each lepton channel is allowed to have its own signal strength is shown in Table 7.5.

Fiducial phase space		Expected $\sigma^{fid,exp}(VVjj)$ [fb]	Measured $\sigma^{fid,obs}(VVjj)$ [fb]
Merged	$\nu\nu qq$	4.1 ± 0.3 (theo.)	10.1 ± 3.3 (stat.) $^{+4.2}_{-3.8}$ (syst.)
	$\ell\nu qq$	6.1 ± 0.5 (theo.)	2.0 ± 1.5 (stat.) $^{+2.9}_{-2.8}$ (syst.)
	$\ell\ell qq$	1.2 ± 0.1 (theo.)	2.4 ± 0.6 (stat.) $^{+0.8}_{-0.7}$ (syst.)
Resolved	$\nu\nu qq$	9.2 ± 0.6 (theo.)	22.8 ± 7.4 (stat.) $^{+9.4}_{-8.5}$ (syst.)
	$\ell\nu qq$	16.4 ± 1.0 (theo.)	5.5 ± 4.1 (stat.) $^{+7.7}_{-7.5}$ (syst.)
	$\ell\ell qq$	6.0 ± 0.4 (theo.)	11.8 ± 3.0 (stat.) $^{+3.8}_{-3.5}$ (syst.)
Inclusive	$\nu\nu qq$	13.3 ± 0.8 (theo.)	32.9 ± 10.7 (stat.) $^{+13.5}_{-12.3}$ (syst.)
	$\ell\nu qq$	22.5 ± 1.5 (theo.)	7.5 ± 5.6 (stat.) $^{+10.5}_{-10.2}$ (syst.)
	$\ell\ell qq$	7.2 ± 0.4 (theo.)	14.2 ± 3.6 (stat.) $^{+4.6}_{-4.2}$ (syst.)

Table 7.5: Expected and measured fiducial cross-sections for the individual lepton channels when allowing the merged and resolved channels to be fit with separate signal strengths or a single common one.

Summary

The electroweak $VVjj$ cross-section in the semileptonic channels was measured in a fiducial region to be 45.1 ± 8.6 (stat.) $^{+15.9}_{-14.6}$ (syst.) fb. The analysis is designed so that the measured $VVjj$ contribution is predominantly due to the vector-boson scattering process. The analysis procedure used 9 different signal regions and 12 control regions for optimal signal acceptance and to accurately measure the backgrounds. This analysis procedure was motivated by the success of the diboson resonance search in the $\ell\nu qq$ channel detailed in Chapter 5 as well as the effectiveness of combined searches discussed in Chapter 6. The development of Boosted Decision Tree classifiers to optimize the signal to background discrimination in this analysis improved the search sensitivity by greater than 70%.

The background only hypothesis p -value is found to correspond to a signal significance of 2.7σ , slightly below the 3σ standard used to claim “evidence” for a signal. The limiting factors were due to both finite data statistics and experimental systematics. The prospects of this analysis in the HL-LHC era with 3000 fb^{-1} are discussed Chapter 8. While direct constraints on non-resonant new physics were not placed, this result establishes the groundwork for subsequent studies to do so. This is the first analysis of this sort published at 13 TeV, covering all semileptonic states, which was not done in the preceding 8 TeV publication [230].

Chapter 8

Semileptonic Diboson Study Prospects at the HL-LHC

Never let the future disturb you. You will meet it, if you have to, with the same weapons of reason which today arm you against the present.

—Marcus Aurelius, *Meditations*

The results shown in the preceding Chapters 5 - 7 used 36 fb^{-1} of 13 TeV pp LHC data as gathered by the ATLAS experiment from 2015 - 2016. During the years of 2024 - 2026, the LHC will be upgraded with several improvements including: 11 T dipole magnets to converge and separate the beams at the IP, crab cavities to rotate the beams before collision, and new 300 m superconducting links to provide current for the magnet systems [232]. This upgraded machine will be known as the High-Luminosity LHC (HL-LHC) and is planned to produce 3000 fb^{-1} of 14 TeV pp collisions by 2035 - 2040 [232]. During this time, the ATLAS experiment will be improved to handle the higher expected instantaneous luminosity with upgrades to all detector systems, in particular to the ID and TDAQ systems [233].

It is useful to consider the future prospects of the previously discussed analyses in Chapters 5 - 7 for various reasons. First, the results of these studies provide motivation for the design of the HL-LHC and help evaluate whether such an accelerator is useful for the scientific community. Secondly, these studies motivate the physics interest for follow-up searches/measurements in these topics with the larger data-sets provided by the HL-LHC.

This chapter will detail prospects for electroweak diboson studies in the $\ell\nu qq$ channel at an upgrade ATLAS detector with 3000 fb^{-1} [11]. These will include extrapolation of the semileptonic diboson resonance results of Chapter 5 as well as the semileptonic vector-boson scattering (VBS) measurements of Chapter 7. Similar VBS prospect studies are presented for the leptonic decay modes in Ref. [234–238]. To the author’s knowledge there has been no similar study done for expected diboson resonance results at the HL-LHC¹. Alongside

¹Except for di-Higgs resonances [239] which differ from the electroweak vector boson resonances discussed in this thesis as they typically probe different BSM models than those discussed here.

prospects in other topics by the ATLAS, CMS, LHCb, and ALICE experiments and theoretical contributions, these studies were gathered in HL-LHC Yellow Reports in Ref. [240–244].

The overall analysis procedure, discussed in Section 8.1, uses MC predictions with parameterizations of the expected upgraded ATLAS detector to estimate the signal and background distributions. The event selection summarized in Section 8.2 attempts to closely match those of the 13 TeV analyses from Chapters 5 - 7, with differences arising from a reduced amount of information stored at the MC level. Estimations of the relative sizes of uncertainties are detailed in Section 8.3, and the results for diboson resonance searches and vector boson scattering are presented in Section 8.4. Most estimates of detector performance and uncertainties used in the chapter are conservative. We will refer to published 13 TeV results of Chapter 5 and Chapter 7 as the nominal analyses in this chapter. For conciseness, only the differences with respect to the nominal analyses will be discussed below.

8.1 Analysis Procedure

The analysis for these prospect studies is based on MC predictions of the relevant processes with the HL-LHC beam conditions (14 TeV center of mass energy and $\langle\mu\rangle \approx 200$), which are then smeared by parameterizations of the expected ATLAS detector performance [245, 246]. The signal samples and the majority of the backgrounds are generated from the same MC configurations described in Section 7.1 and Section 7, now with a center-of-mass energy of 14 TeV. The pile-up conditions are simulated by overlaying inelastic pp Pythia events, following a Poisson distribution with mean $\langle\mu\rangle = 200$.

The only difference in choice of MC generators with respect to the nominal analyses is that Sherpa is no longer used. The V +jets and SM VV processes are instead simulated with MC@NLO, which requires less computation time with respect to Sherpa, and is found to provide better background modeling after applying selection to enhance the VBS signals. No re-weighting of the V +jets distributions as done in Section 7.2.4 is then needed. For these studies, the V +jets events are generated using the MC@NLO v2.3.2 [130] using the NNPDF30NLO set [121], interfaced to Pythia 8 [133] for showering and hadronization. These samples are simulated at NLO for up to one additional parton and up to two additional partons at LO. SM semileptonic diboson production is simulated using POWHEG-Box v2 [132] with the CT10 PDF [125] set and showered/hadronized using Pythia 6 [200], as opposed to Sherpa as well.

Physics level objects are identified by the truth level particle content of the MC generators (see Section 4.1 for truth definition). Electrons and muons are identified directly from the particle label in the truth event. For consistency with the nominal 13 TeV analyses, τ -leptons, will not be directly considered as charged leptons. Both small- R and large- R jets are clustered using the same algorithms as the nominal analyses on all final state particles excluding photons, charged leptons from $W/Z/\tau$ decays, and neutrinos. The missing transverse energy E_T^{miss} is calculated as the negative of the p_T sum of all final state objects within the detector acceptance

of $|\eta| < 4.5$.

To apply the effects of detector performance on the MC predictions, parameterizations of the reconstruction level quantities are calculated in dedicated MC samples with full GEANT4 simulations of the expected upgraded ATLAS layout and full object reconstruction procedure applied [245, 246]. The ratio of fully simulated and reconstructed objects to the MC truth level objects is used to derive multi-dimensional parameterizations in p_T and η . The parameterization include various experimental effects such as identification efficiencies, tagging performance, and fake rates. These parameterizations are applied on truth level MC predictions to simulate the effects of the detector performance.

A parameterization of the $D_2^{\beta=1}$ sub-structure for large- R truth jets is not derived and so the V -tagger used in the nominal analysis can not be directly applied to the prospect studies. To still include the effects of boson tagging on large- R jets in these studies, the V -tagger is simulated by applying a single topology dependent scale-factor to each MC sample. The tagging efficiency scale-factors are calculated from fully-simulated 13 TeV MC samples as the fraction of events with a large- R jet passing the Run-2 V -tagger to the number of events with large- R jets within $|m(J) - m(W/Z)| < 50$ GeV. Separate scale-factors are calculated for each background and signal processes. The efficiency of different signal and background processes to pass the V -tagger is then simulated on the 14 TeV samples by applying the derived scale-factors to events with large- R jets. The mass window cut is enforced on the 14 TeV samples whenever the V -tagger scale-factors are applied to incorporate the shaping effect of the tagger. The mass window of the 50% working point V -tagger, shown in Figure 5.3, is always smaller than the 50 GeV window used in this study. This represents a conservative estimation of the boson-tagging performance in the HL-LHC era.

8.2 Event Selection

The analysis selection used for the prospect studies is designed to follow the diboson selection of Section 5.2 and the VBS selection of Section 7.2 as closely as possible. Requirements on the detector-smearred truth-level object kinematics directly match those placed on the reconstructed level objects of the nominal analysis. This again refers to a conservative scenario where no major analysis improvements are discovered in the upcoming years.

8.2.1 Object Level Differences

The electron and muon isolation requirements are changed such that a lepton is isolated only if the $\sum p_T$ of truth objects within $\Delta R < 10$ GeV/ $p_T(\ell)$ is less than 6% of the lepton p_T . For compatibility with the nominal analyses, the $|\eta(\ell)| < 2.5$ requirement is kept in the prospect studies as opposed to the possible extended range of $|\eta(\ell)| < 4.0$ provided by the upgraded ATLAS geometry. The effect of the reduced η range is expected to be negligible as the leptons from the considered resonant signals and VBS process are expected to be produced centrally.

It is assumed the effects of trigger in the prospect studies can be ignored as it is expected that future lepton and E_T^{miss} triggers will still be fully efficient for the p_T/E_T ranges of interest.

Both analyses reconstruct the hadronic $V \rightarrow q\bar{q}$ as either two small- R jets, or one large- R jet. These are the merged and resolved channels respectively. There is no separation of W -tagged and Z -tagged, or into high-purity or low-purity regions.

8.2.2 VBS Search Differences

With respect to the nominal semileptonic VBS search, two aspects of the analysis have been changed: the lepton channels and the Boosted Decision Tree (BDT) configuration. For simplicity, only the search sensitivity in $\ell\nu qq$ channel is evaluated directly. To extrapolate to all three semileptonic channels, we will assume that all channels provide equal sensitivity.

In the nominal VBS analysis, the final discriminant is a BDT trained on the fully simulated 13 TeV background and signal samples. Due to limited 14 TeV MC statistics, the BDT for the prospect analysis is also trained on the same 13 TeV fully-reconstructed MC samples using the same procedure of Section 7.3, but excluding observables which are not available at the MC truth level, such as the jet width defined in Equation 7.2. The excluded observables have sub-dominant discrimination power with respect to other observables, so the BDT response is not altered significantly. This represents a conservative estimation that no improvement or degradation in the input variable performance will be found with the HL-LHC ATLAS detector.

8.2.3 Kinematic Distributions

Figure 8.1 shows the expected invariant mass distribution, $m(WV)$, of the reconstructed diboson system for the resonance search for gluon-gluon fusion or quark-anti-quark fusion ($ggF/q\bar{q}$), and vector-boson fusion (VBF) produced new particles. The BDT and invariant mass distributions of the merged and resolved regions of the VBS search are displayed in Figure 8.2.

8.3 Systematics

In comparison to the preceding sections where only differences with respect to the nominal analyses were detailed, this section will detail explicitly how the sources of uncertainties were evaluated. Most uncertainties will be ignored so it will be simpler to detail the few ones under consideration.

It is difficult to estimate the expected size of uncertainties at the HL-LHC with respect to current measurements. On-one-hand the majority of the experimental sources of uncertainty are expected to decrease as the large integrated luminosity will provide data-sets 100 times larger than those used currently. On-the-other-hand, the experimental effects of $\langle\mu\rangle \approx 200$ are difficult to characterize until the HL-LHC starts operation. It will be assumed that in the

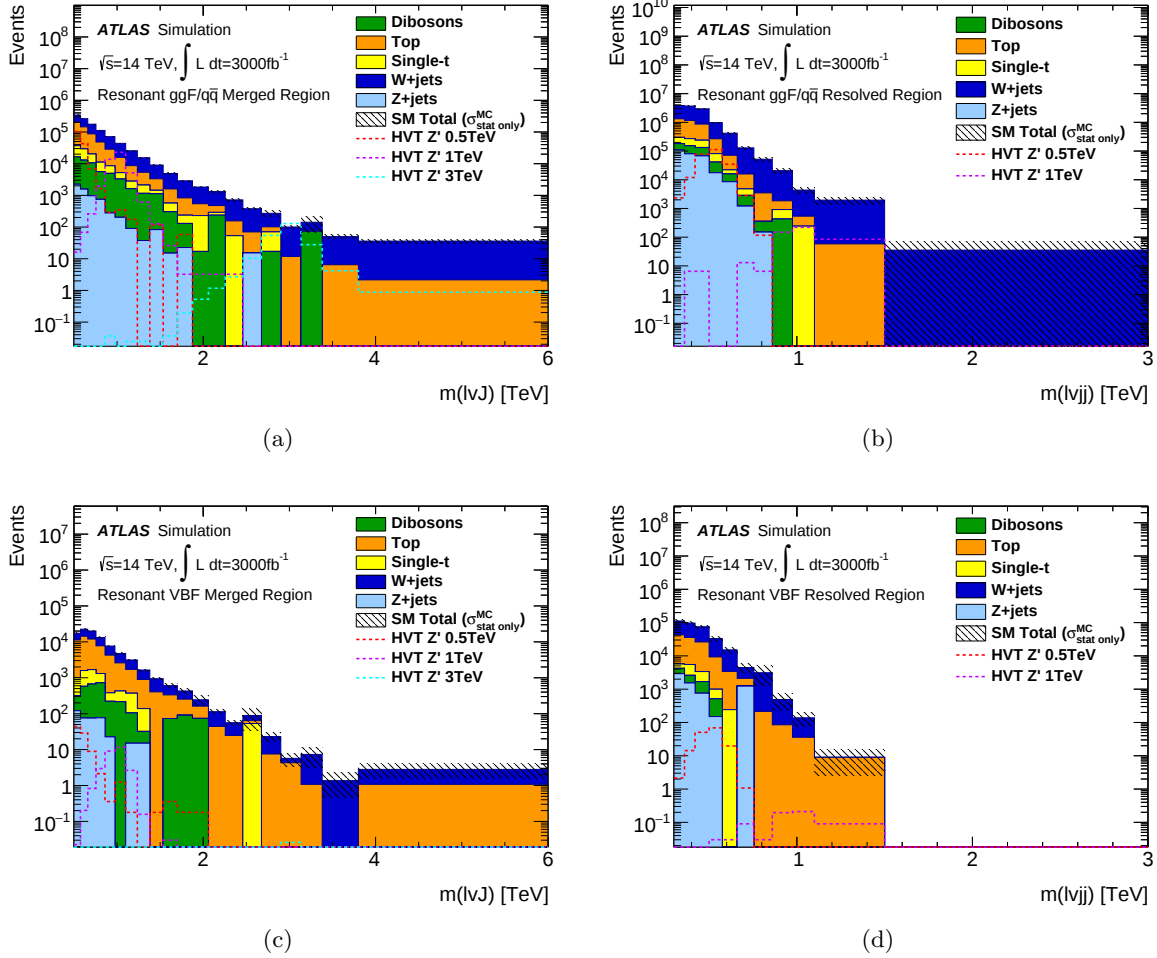


Figure 8.1: Final $m(l\nu jj)$ and $m(l\nu J)$ distributions in the resolved (left) and merged (right) signal regions respectively for the VBF resonance search (top) and ggF/ $q\bar{q}$ resonance search (bottom). Background distributions are separated into production type. HVT signal for masses of 0.5, 1, and 3 TeV are overlaid as dashed curves where appropriate.

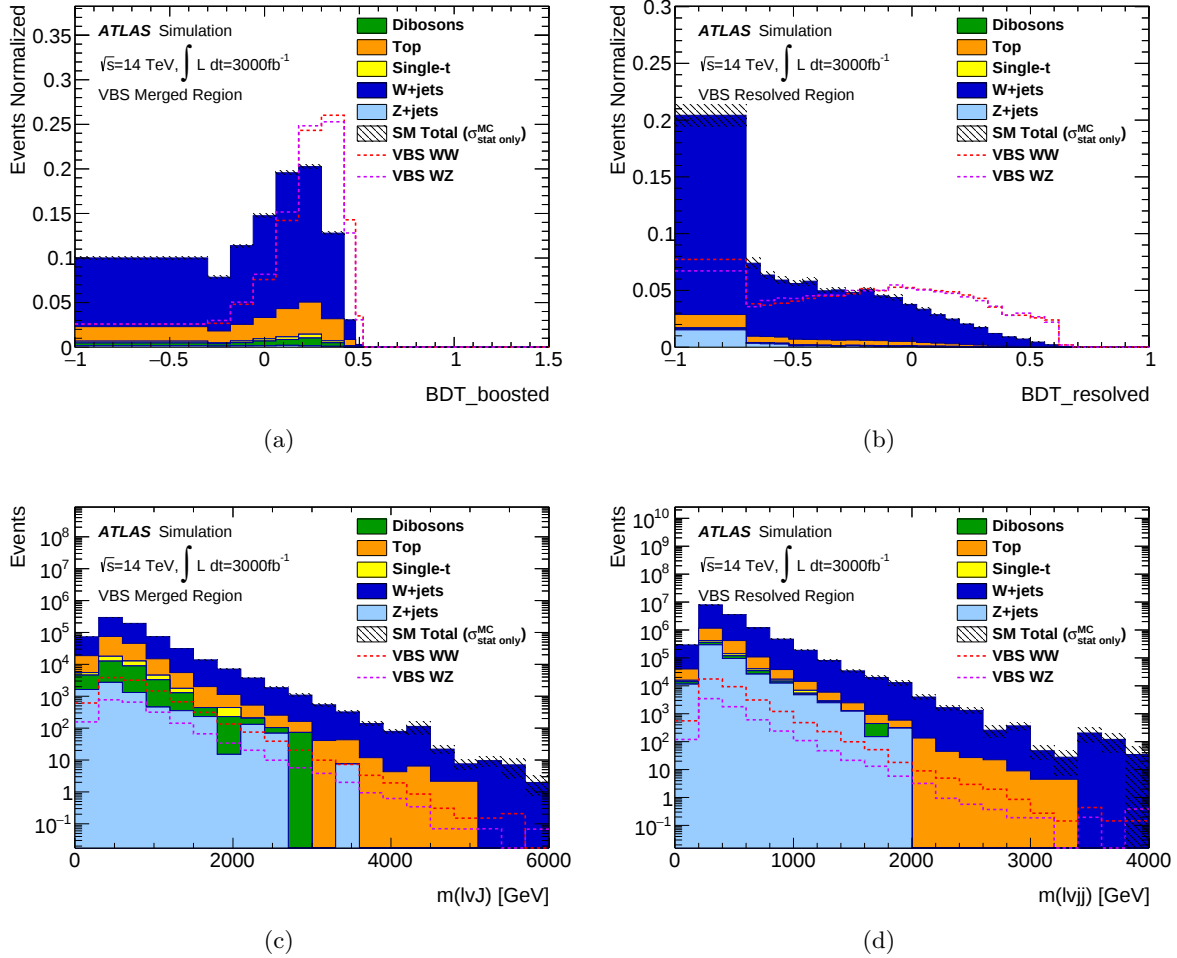


Figure 8.2: Final signal and background distributions for the VBS search in the respective resolved (left) and merged (right) signal regions for the normalized BDT response (top) and the reconstructed diboson invariant mass (bottom). Background distributions are separated into production type. VBS signals in WW and WZ mode are overlaid as dashed curves where appropriate. Both background and signal BDT distributions (top) are normalized to unity

subsequent years of data taking the field will develop improvements in systematic assessment techniques to mitigate these effects.

Only the dominant experimental and theoretical uncertainties from the nominal analyses will be considered here. Other uncertainties are not expected to greatly impact the prospect results. The general procedure will be to apply half the uncertainty as evaluated by the 13 TeV analyses. The uncertainty on the luminosity is set to 1% and the uncertainties due to limited MC statistics are ignored.

The major sources of experimental uncertainty for the resonant and VBS searches in the resolved and merged channels are the small- R and large- R JER respectively. In addition, the large- R JMR uncertainties are considered for the merged channels. The JER and JMR uncertainties used for these results are taken to be half the Run 2 values.

The major background modeling uncertainties on the invariant mass and BDT distributions are those from the W +jets and $t\bar{t}$ backgrounds. For both of these processes, the uncertainty is evaluated using the same procedure as the nominal analysis by a comparison between Madgraph [130] and Sherpa [204] predictions for W +jets, and POWHEG [132] and Madgraph for $t\bar{t}$. The theoretical normalization uncertainties are also considered for all backgrounds and signals. The value of the normalization and shape uncertainties is taken to be half of those from the Run-2 search except for the W +jets which is reduced by a factor of 10. The different scaling of W +jets is chosen due to the expected decrease in cross-section and modeling uncertainties for this process, which is the major background, with a factor of 100 increase in statistics using 3000 fb⁻¹ of data.

8.4 Prospect Results

The final results are extracted from fits to a profile-likelihood model as detailed in Chapter 5 and Chapter 7. Asymptotic formula are used to derive the 95% CL_s upper limits on resonant signal cross-sections as well as the background only hypothesis p -values. The resonant search uses the invariant mass of the $\ell\nu qq$ system as the final observable in the fit while for the VBS search the BDT distribution is used. The binning used for the resonant and VBS studies are taken directly from the nominal analyses.

Benchmarks values are quoted for integrated luminosities of $L = 300$ and $L = 3000$ fb⁻¹, which correspond to the expected data-sets from the LHC Run 3 and HL-LHC operation respectively. The results at $L = 300$ use the same detector configuration and pileup conditions as the HL-LHC studies. Since the analysis strategies are very close to that of the Run 2 nominal analyses and conservative estimations are made, the extrapolation to Run 3 conditions is reasonable.

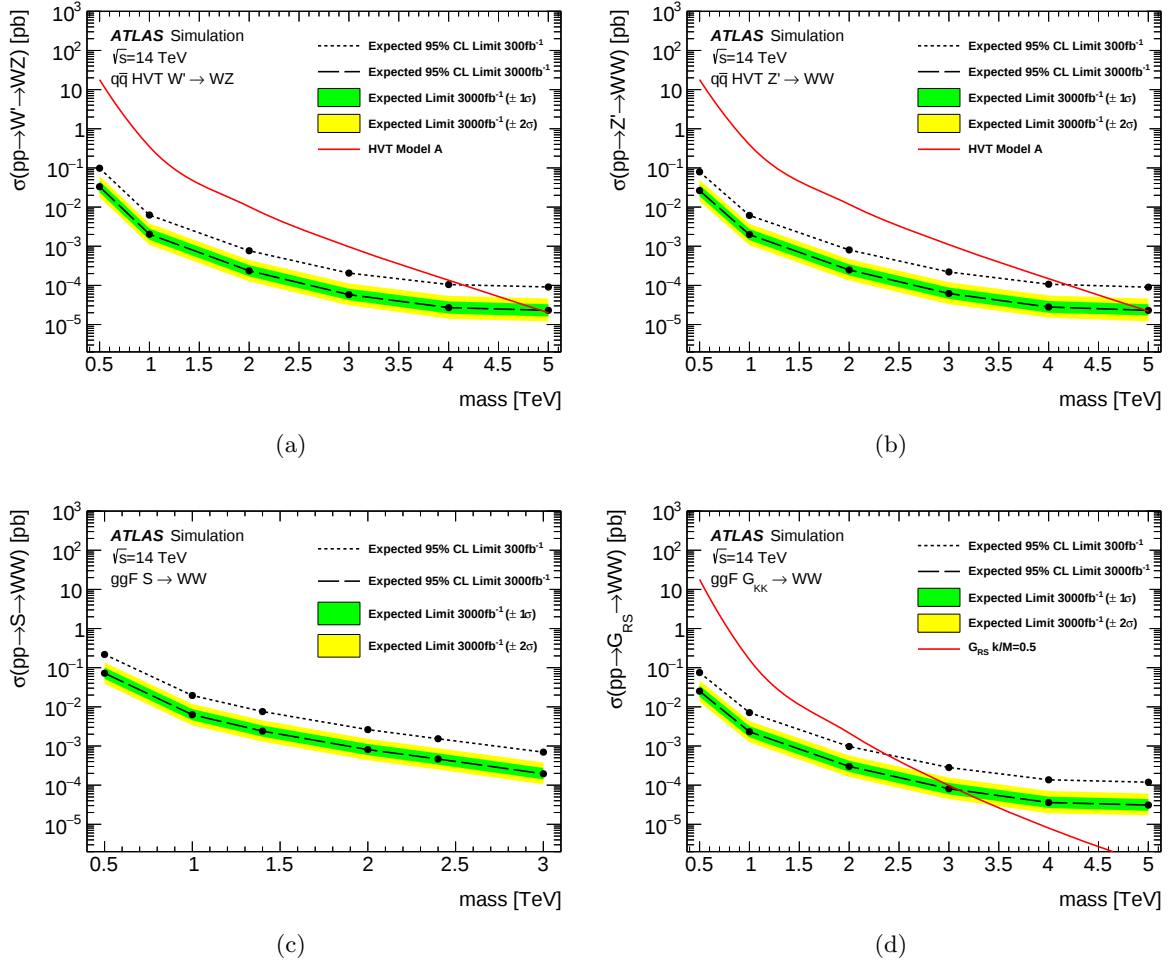


Figure 8.3: The excluded cross-section times branching ratio at 95% CL_s as a function of mass for a) the HVT W' , b) HVT Z' , c) heavy Higgs-like scalar, and d) Bulk RS graviton via $ggF/q\bar{q}$ production with integrated luminosities of 300 (dotted) or 3000 (dashed) fb^{-1} . The expected uncertainties on the 3000 fb^{-1} limit are also shown. The theoretical cross-sections from the HVT Model A and Bulk RS graviton are shown in red.

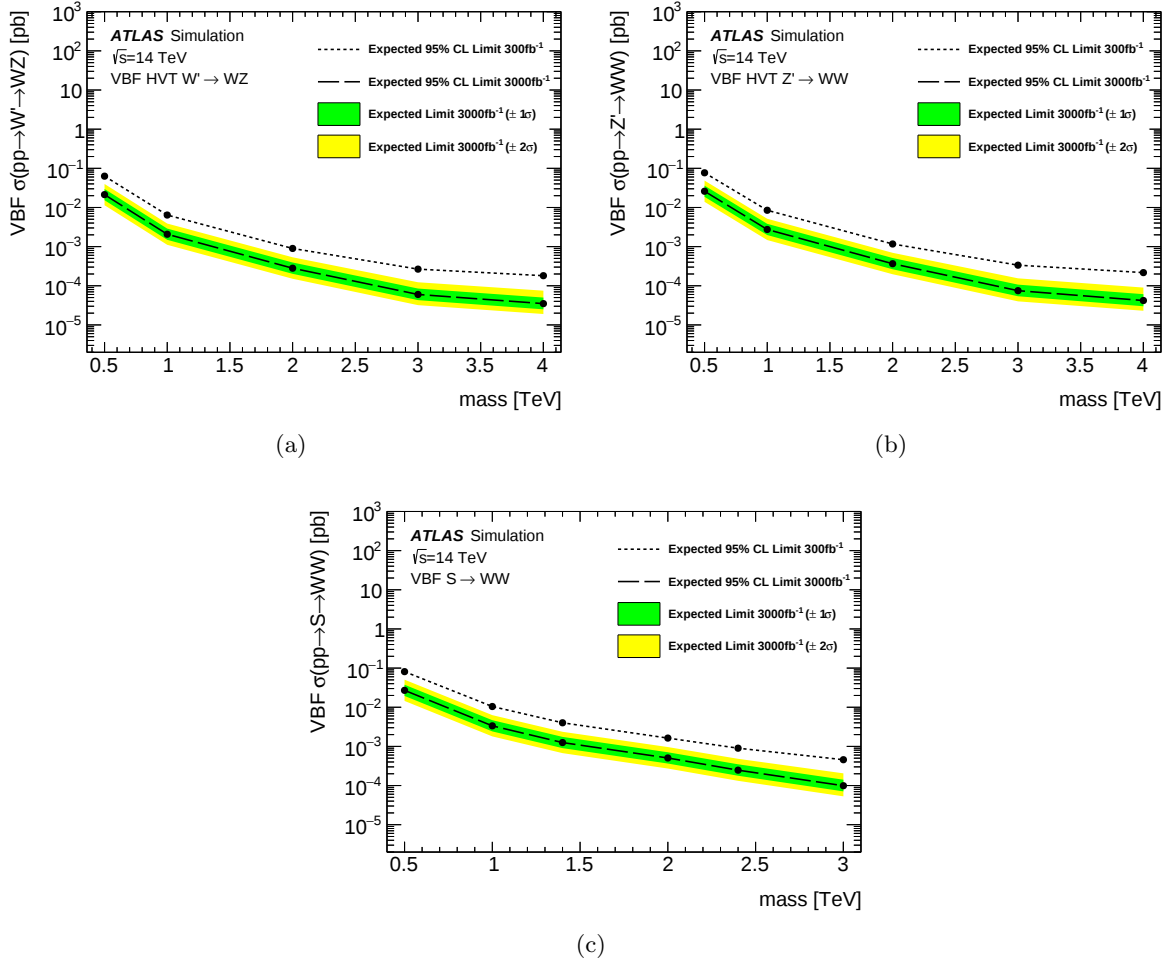


Figure 8.4: The excluded cross-section times branching ratio at 95% CL_s as a function of mass for a) HVT W' , b) HVT Z' (right), and c) heavy Higgs-like scalar via VBF production with integrated luminosities of 300 (dotted) or 3000 (dashed) fb^{-1} . The expected uncertainties on the 3000 fb^{-1} limit are also shown.

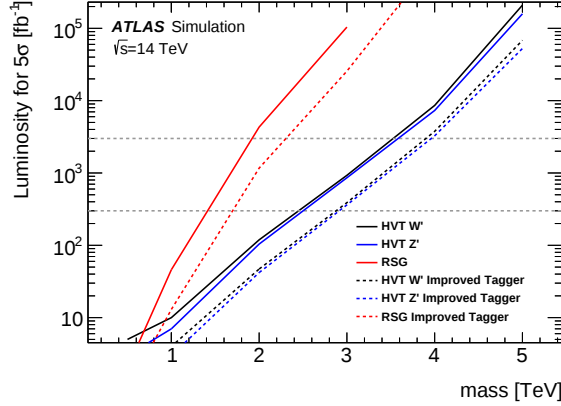


Figure 8.5: Expected luminosity required to claim a 5σ signal significance for the HVT W' (black), HVT Z' (blue), and Bulk RS graviton (red). The solid curves show the sensitivity using the current V -tagger and the dashed curves for a possible future tagger with 50% increased signal efficiency and a factor two increased background rejection.

8.4.1 Resonance Search Prospect Results

The expected 95% CL_s upper limits on the signal cross section times branching ratio as a function of the signal mass in the signal+background hypothesis for different benchmark models are shown in Figures 8.3 - 8.4. Separate projected exclusions are shown for 300 fb^{-1} and 3000 fb^{-1} . The expected upper excluded mass limit for the HVT W' and Z' are estimated to be 4.3 TeV with 300 fb^{-1} and 4.9 TeV with 3000 fb^{-1} . For the Bulk RS graviton, the expected upper mass limits are 2.8 and 3.3 TeV at 300 fb^{-1} and 3000 fb^{-1} . In comparison to the results presented with 36 fb^{-1} in Section 5.4, the 3000 fb^{-1} predictions show improvements on the upper excluded mass by 1.3 TeV for the Bulk RS graviton and 2 TeV for the HVT Model A. This is a relative increase on the upper limits of approximately 75% and 65% respectively.

To study the possible situation where an excess consistent with a diboson resonance signal is observed, the expected discovery luminosity at the HL-LHC is also evaluated. The discovery luminosity is defined as the luminosity required to excluded the background only hypothesis at significance of 5σ in the signal+background hypothesis. For this calculation, the significance is defined as the quadratic sum of $s/\sqrt{s+b}$ for each bin of the final discriminant distribution, where s (b) represent the expected number of signal (background) events in the bin. Figure 8.5 displays the expected discovery luminosity for the resonant search benchmarks as a function of mass. Dashed curves shows the expected values for a potential future V -tagger which has a 50% improvement in the signal efficiency and a factor two improvement in background rejection. These values are representative of initial improvements seen in a diboson resonance search in the fully-hadronic $VV \rightarrow qqqq$ analysis with 80 fb^{-1} of 13 TeV data [247] by using Track-CaloClusters [248] for large- R jets as opposed to LCTopo clusters. Since the publication of these prospect studies it has been found that the initial results of this improvement in the $qqqq$ search of Ref. [247] was incorrect and has since been updated. These values are still representative of

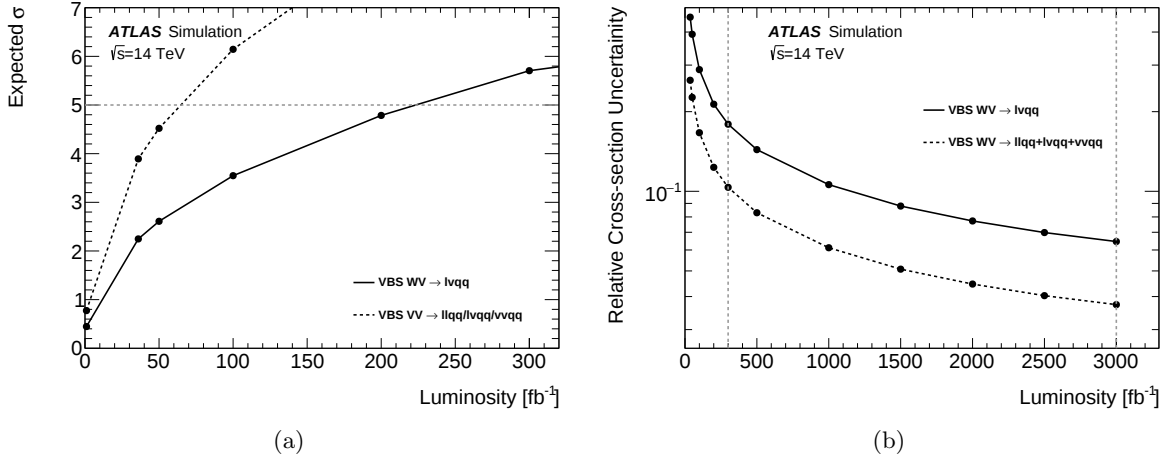


Figure 8.6: a) Expected VBS signal significance as a function of integrated luminosity. The solid black curve is the expected significance from the $\ell\nu qq$ channel, while the black dashed curve shows the expected significance from all semileptonic channels assuming equal sensitivity. The grey dashed curve highlights the 5σ significance value. b) The relative VBS cross-section uncertainty as function of integrated luminosity in the HL-LHC era. The solid black curve is the uncertainty from the $\ell\nu qq$ channel, while the dashed curve shows the expected uncertainty from all semileptonic channels assuming equal sensitivity. The grey dashed curves highlight the values of 300 fb^{-1} and 3000 fb^{-1} . The effects of unfolding are not considered.

reasonable V -tagger improvements which may be developed within the next decade. Possible improvements in V -tagging in the HL-LHC era may originate from more advanced machine-learning techniques to discriminate against the background contribution or better understanding of jet substructure variables with measurements at higher integrated luminosity.

8.4.2 VBS Search Prospects Results

Figure 8.6 shows the expected background only p -value and cross-section uncertainty for VBS in the $\ell\nu qq$ channel as a function of integrated luminosity. Additional curves representing the total combined semileptonic channel (including $\ell\ell qq$ and $\nu\nu qq$) sensitivity are shown assuming each channel has equal sensitivity as the $\ell\nu qq$ channel. The background only p -value for the SM semileptonic VBS process is estimated to be at a significance of 5.7σ at 300 fb^{-1} and the expected cross-section uncertainties are 18% at 300 fb^{-1} and 6.5% at 3000 fb^{-1} . The effects of unfolding detector level observables to truth level predictions, which was implemented using correction factors in Section 7.5.2, were not considered for the cross-section estimates here.

Summary

The studies in this chapter present the prospects of diboson resonant searches and vector-boson scattering measurements in the semileptonic channels in the near future [11]. In the

context of the resonant searches, the expected 95% CL_s limits are expected to increase by greater than 1 TeV with respect to current bounds, dependent on the signal model. In the possible situation where a new diboson resonance could be discovered, it is found in Figure 8.5 that most of the parameter space where we can observe a 5σ excess at 3000 fb^{-1} is excluded by 95% CL_s limits (c.f Figures 6.9 - 6.10). This accentuates the need for analysis improvements in the future such as improved boson tagging or exploiting combinations as discussed in Section 6 to extract the most useful information out of the data.

For the prospects of VBS in the semileptonic channels, it is expected that the background only hypothesis can be excluded at 5σ observation with the full Run 2 data-set and all semileptonic channels. In the HL-LHC era with 3000 fb^{-1} the cross-section of this process is expected to be measured at the level of 4%.

Chapter 9

Summary

Leave her, Johnny, leave her!
Oh, leave her, Johnny, leave her!
For the voyage is done and the winds don't blow
And it's time for us to leave her.
—Historical Sea Shanty, *Leave Her Johnny*

This thesis presents my work on studies in final states with pairs of electroweak vector-bosons decaying into the semileptonic final state. The diboson states provide sensitivity to possible BSM physics or general alteration of the boson couplings. These results utilized 36 fb^{-1} of 13 TeV pp collisions gathered with the ATLAS detector from 2015 - 2016. Each analysis result is separately summarized below.

$VV \rightarrow \ell\nu qq$ Resonance Search

The search for new resonant states which decay to dibosons in the $\ell\nu qq$ final state was performed [8]. The analysis utilizes modern advances in jet substructure to identify jets originating from the merged decays of $V \rightarrow q\bar{q}$, as well as a newly optimized selection to search for resolved $V \rightarrow q\bar{q}$ jets. In comparison to previous studies, a dedicated search was also performed for resonances produced via vector-boson-fusion (VBF).

All signal regions were found to be consistent with the background only prediction. Upper limits were set on the possible resonant mass in the context of several benchmark models of varying spin. The upper mass limits at 95% CL_s for the Heavy Vector Triplet (HVT) were found to be at 2.8 TeV for a benchmark coupling choice reminiscent of a weakly coupled model (Model A). For a strongly coupled model choice (Model B) the limits were found to be at 3.0 TeV. Similar limits were set on a bulk Randall-Sundrum Graviton with $k/M_{Pl} = 1$ at 1.7 TeV.

Diboson+Dilepton Resonance Combination

The above quoted diboson limits were only provided for a single decay channel to the possible resonance. If new resonances were to exist, one would expect the results of all searches in each diboson decay channel to be correlated. In general, a new particle could decay into multiple intermediate/final states, so several analyses together can provide increased constraining powers on such a model. Results were presented on the combination of all the previously published ATLAS results for new physics in the VV , Vh , and dilepton channels [9].

The data was again found to be consistent with the background only prediction after combining results from 14 separate publications. Limits were shown using possible model independent combinations of the VV , Vh , and $\ell\ell$ channels independently, followed by model dependent total combination in the context of the HVT model. The VV combination increased the observed upper mass limits in the HVT Model A to 2.8 TeV and the $VV+Vh+\ell\ell$ combination to 5.5 TeV. Direct limits on the possible resonance couplings to SM particles in the HVT model were also presented. These improved upon existing constraints from electroweak precision tests.

Semileptonic VBS search

A search for the SM production of electroweak vector boson scattering (VBS) was also performed in the semileptonic channels [191]. The main analysis strategy was taken from the VBF $\ell\nu qq$ diboson resonance search and re-optimized for a SM measurement. This result also utilized the other semileptonic channels: $\ell\ell qq$ and $\nu\nu qq$. To optimize the search sensitivity a Boosted Decision Tree Classifier was implemented.

Disagreement between the background only assumption without VBS was found at the level of 2.7σ . The best fit value to data of the observed signal strength relative to the SM prediction is $\mu = 1.05^{+0.42}_{-0.40}$. A measurement of the cross-section was performed in a fiducial volume for direct comparison to theoretical predictions. The fiducial semileptonic VBS cross-section was measured to be

$$\sigma = 45.1 \pm 8.6 \text{ (stat.)}^{+15.9}_{-14.6} \text{ (syst.) fb}$$

Semileptonic Diboson Study Prospects at the HL-LHC

The prospects of the above studies in the HL-LHC era were also evaluated [11]. This was done using MC predictions and simulations of the upgraded ATLAS geometry plus performance. With 3000 fb^{-1} of data, the upper excluded mass limits for $\ell\nu qq$ resonance search limits are 4.9 TeV for the HVT W'/Z' in the benchmark Model A and 3.3 TeV for the Bulk RS graviton. This is an increase by approximately 2 TeV for the HVT Model A and 1.3 TeV for the Bulk RS graviton benchmark with respect to current 36 fb^{-1} limits. For semileptonic VBS, a discovery significance greater than 5σ is expected with the full Run-2 data and with 3000 fb^{-1} the cross-section is predicted to be measurable at the percent level.

Appendices

Appendix A

Statistical Framework

It is frequent for experiments in high-energy physics to measure the number of events as a function of an observable. The expected distribution of the signal and background are then fit to the observed data distribution. The results of this fit can be interpreted in different ways depending on the context of the analysis. If an analysis is searching for a signal, the compatibility of the background only fit to data is quantified by a p -value. For an observed signal, the fit results can be interpreted as a measurement of the signal cross-section and its uncertainty. In other circumstances, the results can be used to measure which signal model parameters are not excluded by the data.

In this section, the statistical techniques used to evaluate the interval of hypothesis parameters excluded by the data and background only p -values are discussed in detail. Section A.1 describes the likelihood model used for the fit as well as the test statistic used to evaluate the compatibility of data and the fit. The test statistic is based on the profile-likelihood ratio of the data fit to different hypothesis. The derivation of a simplified asymptotic formula for evaluation of these test statistics in the large sample limit is given in Section A.2.

A.1 Profile-likelihoods

The analysis follows a frequentist approach, where a likelihood model is built and a test statistic based on the profile-likelihood ratio is used for hypothesis testing [249–251].

We will assume that the analysis will measure the number of data events in bins of some observable x , and that the expected signal and background yields can be estimated in these bins. Explicitly for Chapter 5 this is the diboson invariant mass $m(VV)$ and for Chapter 7 this is the BDT score, with both signal and background distributions estimated from MC. The signal cross-section will be parameterized by a scale-factor μ . The value of $\mu = 0$ then corresponds to the background only prediction, and $\mu = 1$ the signal plus background prediction. The final parameter of interest will be μ , which will frequently be allowed to float in fits to the best value which represents the data.

The prototype for our likelihood model is that of a simple binned counting experiment. In

such a case the likelihood is the product of Poisson probabilities to measure n_i events with s_i signal and b_i background events expected

$$L(\mu) = \prod_i Pois(n_i | \mu s_i + b_i) = \prod_i e^{-(\mu s_i + b_i)} \frac{(\mu s_i + b_i)^{n_i}}{n_i!} \quad (\text{A.1})$$

where i indexes the product over all the bins. It is typical for analyses, especially those which rely on MC estimates, to derive template functions of the expected background $f_b(x)$ and signal $f_s(x)$ distributions as a function of the observable x . The expected binned distribution yields can be extracted from the templates by

$$s_i = \int_{bin\ i} f_s(x) dx \quad b_i = \int_{bin\ i} b_s(x) dx \quad (\text{A.2})$$

In this analysis we expand the above definition of the likelihood to include the effects of systematics as additional nuisance parameters¹ θ_k . The template functions are now parameterized by these systematics, $f_s(x, \boldsymbol{\theta})$ and $f_b(x, \boldsymbol{\theta})$, which encode the possible shape and normalization effect associated with each of the sources of uncertainty.

The nuisance parameters are constrained by additional terms in the likelihood $f(\theta_k | \boldsymbol{\alpha}_k)$, which are dependent on external measurements α_k which parameterize the expected probability distribution of the θ_k measurement. A simple illustration of this for example, is if θ_k represented the small- R JER uncertainty, then the constraint function could be a Gaussian where $\boldsymbol{\alpha}_k$ contain the externally measured central value and standard deviation of the JER. The statistical uncertainties related from finite MC statistics can also be included using a similar method with Poisson constraint [252].

The likelihood function including the additional systematic terms is then

$$L(\mu, \boldsymbol{\theta}) = \prod_i Pois(n_i | \mu s_i(\boldsymbol{\theta}) + b_i(\boldsymbol{\theta})) \prod_k f(\theta_k | \boldsymbol{\alpha}_k) \quad (\text{A.3})$$

The combination of several regions, either new signal or control regions, is simply the addition of new bins in the product i , with careful consideration of which θ_k are common. A “fit” to a data-set is the substitution of the maximum likelihood estimates of the model parameters as evaluated from that data-set. These maximum-likelihood values serve as point estimates, and the interval estimates can be conducted in two ways. The first method, known in particle physics as the “Hess” method, estimates the confidence interval of these point by estimating the covariance matrix around the likelihood minimum, and using a parabolic expansion around this result. In the second method, sometimes known as the “Minos” method, the confidence interval is estimated by the domain in which the likelihood differs by less than a threshold value with respect to the minimum. For a single parameter and a 68% confidence interval, this would

¹In this chapter we will denote a vector of parameters with a bold notation (e.x. $\boldsymbol{\theta} = (\theta_1, \theta_2, \dots, \theta_n)$ or $\boldsymbol{\alpha}_k = (\alpha_{k,1}, \alpha_{k,2}, \dots, \alpha_{k,n})$).

be when the likelihood differs by 1/2 with respect to the minimum.

The final test statistic is based on the profile likelihood ratio [249]

$$\lambda(\mu) = \frac{L(\mu|\hat{\boldsymbol{\theta}}(\mu))}{L(\hat{\mu}|\hat{\boldsymbol{\theta}})} \quad (\text{A.4})$$

where the single hat notation $\hat{\boldsymbol{\theta}}$ indicates the maximum likelihood fit value, and $\hat{\boldsymbol{\theta}}(\mu)$ indicates the conditional maximum likelihood value for a specific choice of μ . The motivation for using the profile likelihood ratio is the Neyman-Pearson lemma which states that for single-parameter cases, the likelihood ratio forms the most powerful test statistic for a given significance level [253]. While the Neyman-Pearson lemma does not hold for multi-dimensional likelihoods such as Equation A.3, we attempt to approach such a situation by replacing the nuisance parameters by their profiled values. In a sense, the nuisance parameters are profiled away, and the effective dimensionality of the likelihood reduced. The test statistic $0 < \lambda(\mu) < 1$ measures the level of agreement to data for a specific value of μ , with $\lambda(\mu) = 1$ indicating the best agreement.

In the context of searching for new physics, $\hat{\mu} < 0$ indicates there is a deficit of events in data with respect to the background prediction. Since we are searching for solely positive excesses, we define the test statistics such that for $\hat{\mu} < 0$ the best fit value should be for $\mu = 0$. To quantify an excess the following test statistic is thus used

$$q_0 = \begin{cases} -2 \ln \frac{L(0|\hat{\boldsymbol{\theta}}(\mu))}{L(\hat{\mu}|\hat{\boldsymbol{\theta}})} & \hat{\mu} > 0, \\ 0 & \hat{\mu} < 0 \end{cases} \quad (\text{A.5})$$

Higher value of q_0 indicate a worse level of agreement between the background only prediction and the best fit (positive) μ value.

To quantify the disagreement between a model and data the test statistic is

$$\tilde{q}_\mu = \begin{cases} -2 \ln \frac{L(\mu|\hat{\boldsymbol{\theta}}(\mu))}{L(0|\hat{\boldsymbol{\theta}}(0))} & \hat{\mu} < 0, \\ -2 \ln \frac{L(\mu|\hat{\boldsymbol{\theta}}(\mu))}{L(\hat{\mu}|\hat{\boldsymbol{\theta}})} & 0 < \hat{\mu} < \mu \\ 0 & \mu < \hat{\mu} \end{cases} \quad (\text{A.6})$$

Again, higher values of $\tilde{q}_\mu$ indicate a worse agreement of the data with the signal of strength μ . When $\hat{\mu} > \mu$ we still view the data as being consistent with that μ value.

The p -values for a one-sided test can be extracted by integrating the expected probability distribution of the test-statistic above the measured test statistic value. We will often convert p -values to a “signal significance” defined by the transform $n = \Phi^{-1}(1 - p)$, where Φ^{-1} is the inverse of the Gaussian cumulative distribution function. A test with a p -value corresponding to a signal significance n , will be denoted as having a significance of $n\sigma$.

When quoting confidence intervals on excluded values, we will also make use of the CL_s

technique [254]. A signal model is excluded by CL_s test of significance α when

$$CL_s = \frac{p_{s+b}}{1 - p_b} < \alpha \quad (\text{A.7})$$

where p_{s+b} is the p -value of the signal plus background hypothesis, and p_b is the background only hypothesis. This procedure protect against excluding models in searches where one has almost no sensitivity by adjusting the normal p -value by the factor of $1 - p_b$. For a test with p_b approximately zero, the CL_s method matches the typical p -value.

A.2 Asymptotic Formula for Test Statistic

The evaluation of the test statistic distributions $f(\tilde{q}_\mu|\mu')$ and $f(q_0|\mu')$ for a given μ' can be computationally expensive using normal Monte-Carlo methods. The task can be simplified through use of analytic asymptotic formula valid in the large sample limit [249]. These formulae utilize results from Wilks [255] and Wald [256], which shows that the $-2\log(\lambda(\mu))$ distribution asymptotically approaches a non-central chi-square distribution. More precisely, in the large sample limit the random variable

$$-2\log\lambda(\mu) \approx \frac{(\mu - \hat{\mu})^2}{\sigma^2} \quad (\text{A.8})$$

where $\hat{\mu}$ follows a Gaussian distribution around the tested value μ' with variance σ . An important technique to estimate the variance is through use of an ‘‘Asimov’’ data-set designed such that expectation value of all the parameters evaluated on this data-set are the true parameters. This is true when the data matches identically to the predictions, $n_i = \mu' s_i(\boldsymbol{\theta}) + b_i(\boldsymbol{\theta})$. The variance can then be estimated quickly as

$$\sigma^2 \approx \frac{(\mu - \mu')^2}{-2\ln\lambda_A(\mu)} \quad (\text{A.9})$$

where $-2\ln\lambda_A(\mu)$ is the profile likelihood ratio evaluated on the Asimov data-set.

For use in limit setting, the probability distribution of the test statistic $f(\tilde{q}_\mu|\mu')$, and the cumulative distributions $F(\tilde{q}_\mu|\mu')$, in the asymptotic approximation can be evaluated as

$$f(\tilde{q}_\mu|\mu') = \Phi\left(\frac{\mu' - \mu}{\sigma}\right)\delta(\tilde{q}_\mu) + \begin{cases} \frac{1}{2\sqrt{2\pi}\sqrt{\tilde{q}_\mu}} \exp\left(-\frac{1}{2}\left(\sqrt{\tilde{q}_\mu} - \frac{\mu - \mu'}{\sigma}\right)^2\right) & 0 < \tilde{q}_\mu \leq \mu^2/\sigma^2 \\ \frac{1}{\sqrt{2\pi}(2\mu/\sigma)} \exp\left(-\frac{1}{2}\frac{(\tilde{q}_\mu - (\mu^2 - 2\mu\mu')/\sigma^2)^2}{(2\mu/\sigma)^2}\right) & \mu^2/\sigma^2 < \tilde{q}_\mu \end{cases} \quad (\text{A.10})$$

$$F(\tilde{q}_\mu|\mu') = \begin{cases} \Phi\left(q_\mu - \frac{\mu - \mu'}{\sigma}\right) & 0 < \tilde{q}_\mu \leq \mu^2/\sigma^2 \\ \Phi\left(\frac{\tilde{q}_\mu - (\mu^2 - 2\mu\mu')/\sigma^2}{2\mu/\sigma}\right) & \mu^2/\sigma^2 < \tilde{q}_\mu \end{cases} \quad (\text{A.11})$$

The upper limit on μ at significance α can then be calculated by inverting $1 - F(\tilde{q}_\mu|\mu) = \alpha$.

The expected limits can be calculated by applying the same procedure to the Asimov data-set for $\mu' = 0$. This leads to a median upper limit value given by $\mu' + \sigma\Phi(1 - \alpha)$ and $\pm N\sigma$ error bands given by $\mu' + \sigma(\Phi(1 - \alpha) \pm N)$.

For background only tests, the probability and cumulative distribution of q_0 given μ' can also be approximated in the large sample limit as

$$f(q_0|\mu') = (1 - \Phi(\frac{\mu'}{\sigma}))\delta(q_0) + \frac{1}{2} \frac{1}{\sqrt{2\pi}\sqrt{q_0}} \exp\left(\frac{1}{2}(\sqrt{q_0} - \frac{\mu'}{\sigma})^2\right) \quad (\text{A.12})$$

$$F(q_0|\mu') = \Phi(\sqrt{q_0} - \frac{\mu'}{\sigma}) \quad (\text{A.13})$$

The local background only p -value can be evaluated as $p = 1 - F(q_0|0)$, and the expected p -value can be determined using an Asimov data-set generated with $\mu' = 0$.

Appendix B

Additional Systematic and Statistical Studies

This appendix includes some supporting details on fit studies in the $\ell\nu qq$ resonance search and the diboson and dilepton resonance combination search of Chapter 6

B.1 Fit Studies

This section will include the sanity checks of the fits as evaluated on the signal plus background templates of the searches. These checks are important for understanding the validity of the fit model and systematic evaluations, but are typically not shown in ATLAS publications. Fit studies were in particular important for the exotic resonance combination of Chapter 6, where care had to be taken that the fit was robust-enough to handle all of the 14 input channels and that all systematics were correlated.

Two sets of plots will be shown. The first set of plots, known as pull plots, show the difference between the fitted value of nuisance parameter θ_{fit} and the externally measured value θ_0 . The uncertainty on θ_{fit} is measured from the curvature at the minimum. Pull plots are expected to be centered at zero with the uncertainty similar to the externally measured uncertainty $\Delta\theta$. Nuisance parameters with smaller uncertainties (constraints) indicate the analysis is possibly a more sensitive measure of the uncertainty than the external measurement. Larger uncertainties (underconstraints) indicate issues with the fits, as well as value of $\frac{|\theta_{fit}-\theta_0|}{\Delta\theta} > 1$ (pulls).

The second set of plots show the linear correlation coefficient of two nuisance parameters X and Y as estimated from the covariance matrix σ_{XY}^2 as $\rho = \frac{\sigma_{XY}^2}{\sigma_X\sigma_Y}$, where σ_X is the root of the diagonal value of the covariance matrix (the standard deviation of X). The covariance matrix is estimated from the curvature of the minimum of the likelihood. Large correlations between nuisance parameters should be understood as they indicate degeneracy within the fit model.

In the subsequent pull and correlation plots, the labels of the systematics in each bin will be suppressed, due to the large number of considered systematics. We will find that most

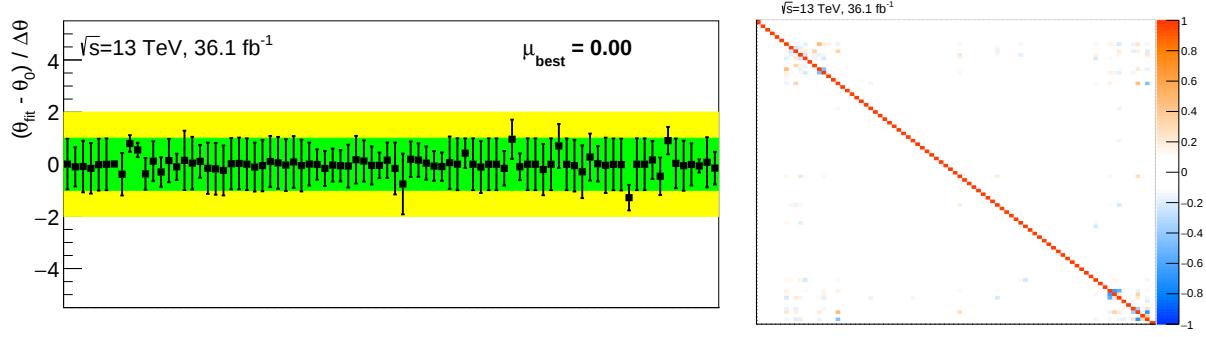


Figure B.1: Fit results from the $\ell\nu qq$ resonance search for a 1.6 TeV $ggF/q\bar{q} WW$ resonance. Left) The relative difference between the fit value of nuisance parameters with respect to external measurement. Right) The correlation coefficient between nuisance parameters.

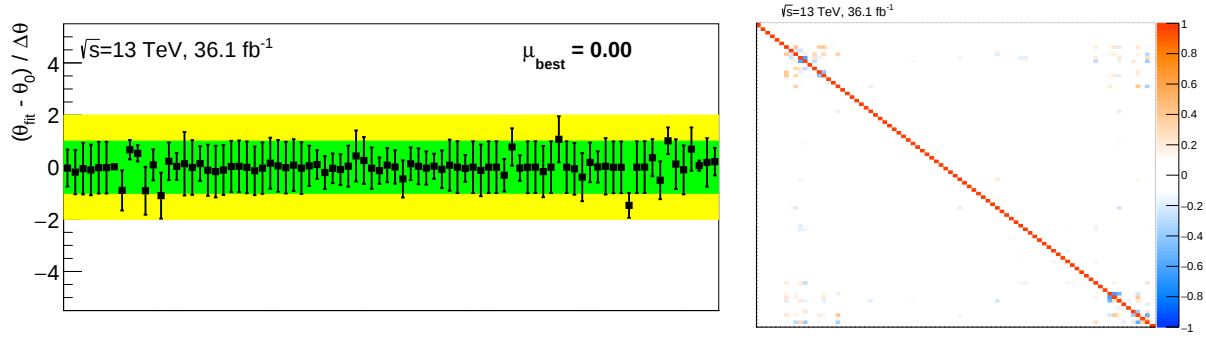


Figure B.2: Fit results from the $\ell\nu qq$ resonance search for a 1.6 TeV $ggF/q\bar{q} WZ$ resonance. Left) The relative difference between the fit value of nuisance parameters with respect to external measurement. Right) The correlation coefficient between nuisance parameters. Each bin corresponds to a separate systematic, with the labels suppressed for readability.

systematics are well behaved in the fit, and any features will be discussed explicitly in the text.

B.1.1 $\ell\nu qq$ Analysis

The fit studies for the $\ell\nu qq$ resonance search of the Chapter 5 in the $ggF/q\bar{q}$ region are shown in Figures B.1 - B.2 and for the VBF search in Figures B.3 - B.4. The fits look well-behaved in all regions. The systematics which are found to be constrained in the fit are associated to the large- R jet uncertainties. These constraints are expected since they are estimated from 20 fb^{-1} of 8 TeV data, while this analysis used 36 fb^{-1} of 13 TeV. Since this analysis is sensitive to the large- R calibration and uses a larger data-set, constraints are expected. Several small- R jet systematics are also found to be slightly constrained for similar reasons. All other pulls and constraints are understood as expected statistical fluctuations. The correlation matrix is found to have only small off-diagonal terms. The small off-diagonal entries are found to originate from large- R jet systematics, which is expected, and the W +jets modeling uncertainties.

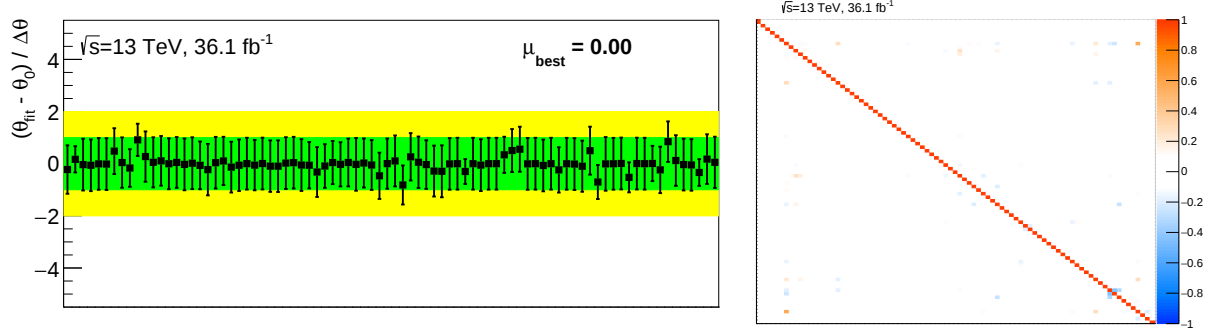


Figure B.3: Fit results from the $\ell\nu q\bar{q}$ resonance search for a 1.6 TeV VBF WW resonance. Left) The relative difference between the fit value of nuisance parameters with respect to external measurement. Right) The correlation coefficient between nuisance parameters. Each bin corresponds to a separate systematic, with the labels suppressed for readability.

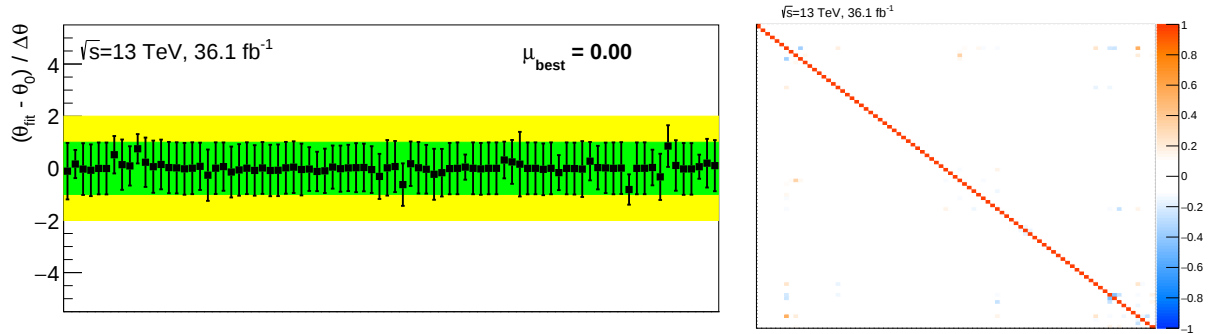


Figure B.4: Fit results from the $\ell\nu q\bar{q}$ resonance search for a 1.6 TeV VBF WZ resonance. Left) The relative difference between the fit value of nuisance parameters with respect to external measurement. Right) The correlation coefficient between nuisance parameters. Each bin corresponds to a separate systematic, with the labels suppressed for readability.

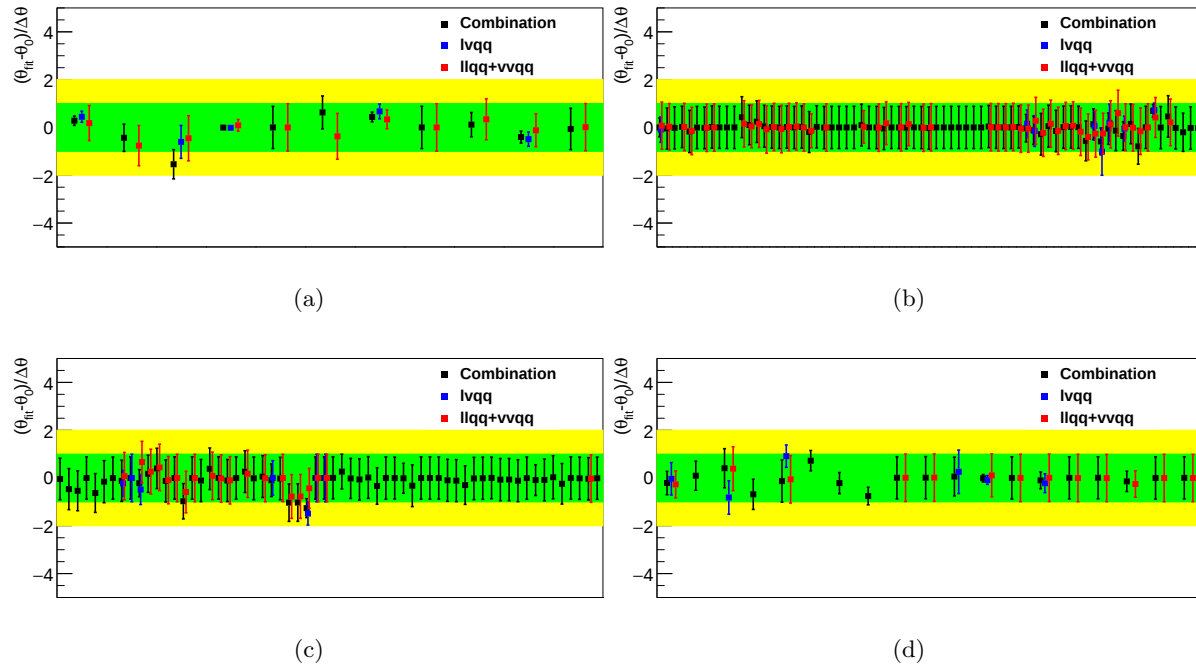


Figure B.5: Pulls from the HVT $WW+WZ$ diboson combination including all relevant channels (black) and from just the $\ell\nu qq$ (blue) and $\ell\ell qq + \nu\nu qq$ channels (red). a) The large- R jet systematics, b) small- R jet systematics, c) other experimental systematics, and d) the theory modeling uncertainties are shown separately.

B.1.2 Diboson and Dilepton Combination

The combination of diboson and dilepton resonance searches in Chapter 6 contains over 100 nuisance parameters. Many systematic studies were done to validate the stability of the fit and the correlation scheme. A comparison of the pulls between the global $WW + WZ$ fit and $WW + ZZ$ with respect to the individual $\ell\nu qq$ and $\ell\ell qq + \nu\nu qq$ fits are shown in Figure B.5 and Figure B.6. Almost all of the pulled and constrained nuisance parameters originate from the $\ell\nu qq$ channel or the $\ell\ell qq + \nu\nu qq$ channel results. The origin of these constraints are understood in the individual publication results. Some additional pulls/constraints are found to originate from the $qqqq$ channel. The total $VV + Vh + \ell\ell$ pulls and correlations are identical to VV . Overall the Vh fits are found to behave similarly as the VV due to similarities in the analysis strategies and final state. The $\ell\ell$ are sensitive only to the leptonic nuisance parameters, which are not constrained in the $VV+Vh$ fits and so do not greatly influence the fit stability.

B.2 Local p -values

Alongside exclusion limits, studies on the local background-only p -value for each of the resonance searches were also conducted. The local p -value was calculated using the asymptotic formula of Equation A.6. The p -values are local in the sense they do not account for the

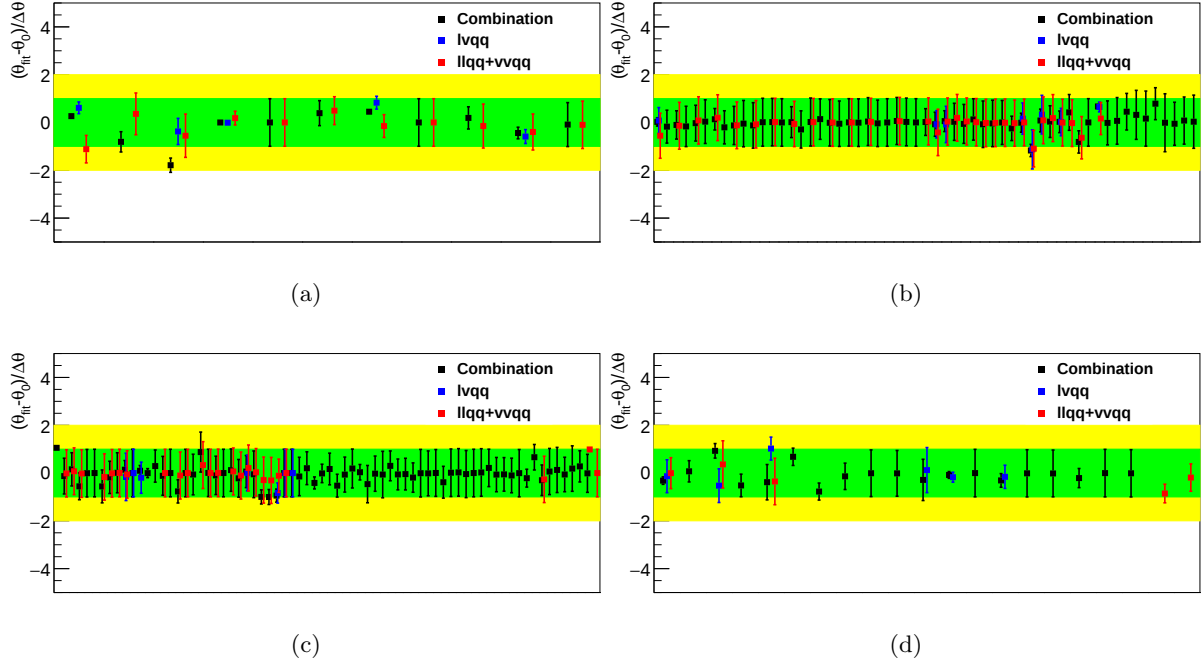


Figure B.6: Pulls from the Bulk RS graviton $WW + ZZ$ diboson combination including all relevant channels (black) and from just the $\ell\nu qq$ (blue) and $\ell\ell qq + \nu\nu qq$ channels (red). a) The large- R jet systematics, b) small- R jet systematics, c) other experimental systematics, and d) the theory modeling uncertainties are shown separately.

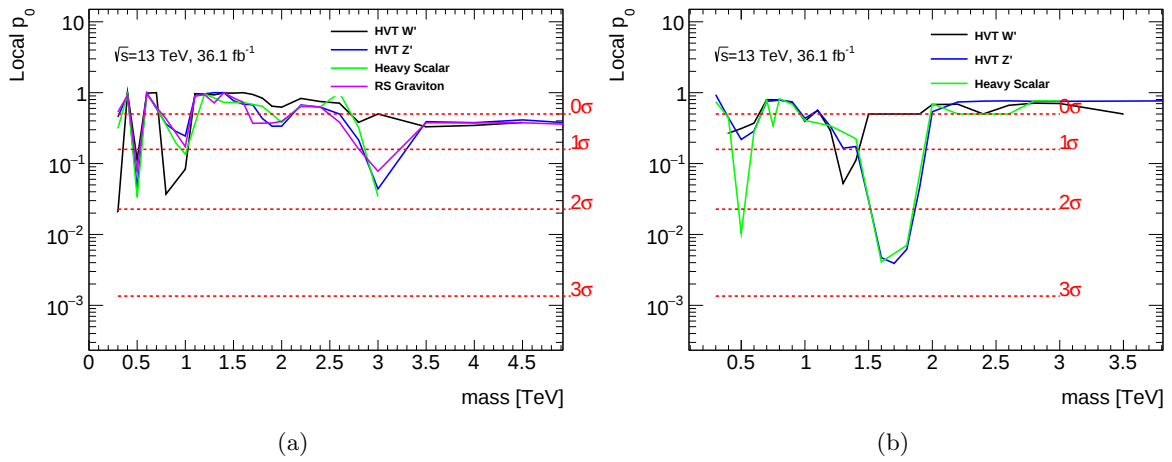


Figure B.7: The local background-only p -value of the a) $ggF/q\bar{q}$ and b) VBF $\ell\nu qq$ resonance search as a function of resonance mass for the different benchmark models.

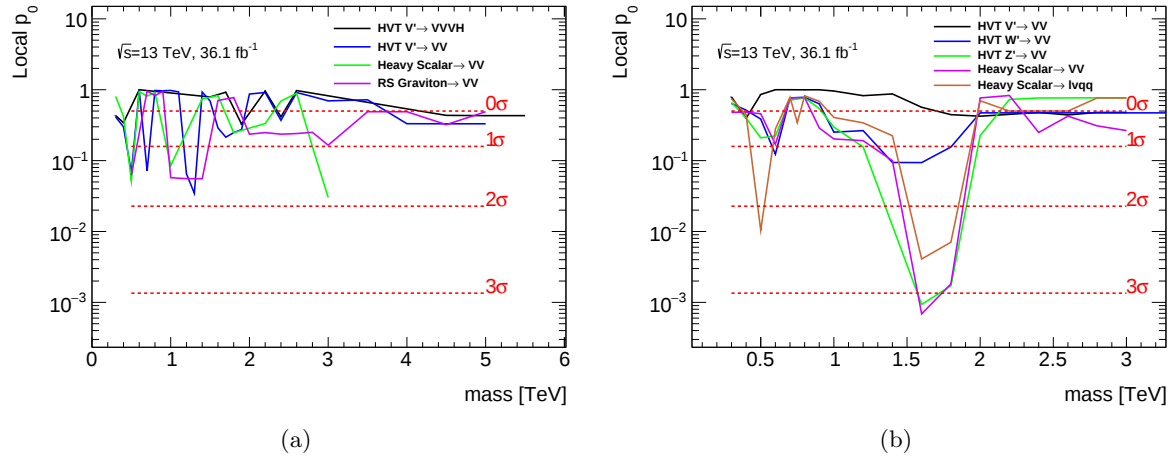


Figure B.8: The local background-only p -value of the a) $ggF/q\bar{q}$ and b) VBF diboson combination resonance search as a function of resonance mass for the different benchmark models.

“look-elsewhere” effects [257].

The background-only p -values for the $\ell\nu q\bar{q}$ resonance search are seen in Figure B.7. The largest discrepancy with the background prediction is 2.7σ in the VBF WW channels. This is consistent with the small excess found approximately at 1.6 TeV in the VBF WW data/MC comparison in Figure 5.6.

The same set of plots for the $VV+Vh$ combinations are shown in Figure B.8. For simplicity in the diboson combination, the $\ell\nu q\bar{q}$ channel had several nuisance parameters pruned, which slightly raises the p -value to above 3σ . This should be identified as an effect of the pruning and not of genuine “evidence”, for new physics.

B.3 Systematics included in the Diboson+Dilepton Resonance Search

This section describes the sources of systematic uncertainty considered in the diboson and dilepton combination detailed in Chapter 6. Almost all of these sources have been previously discussed in the context of the $\ell\nu q\bar{q}$ resonance search in Section 5.3. No analysis considers additional experimental uncertainties than that of the $\ell\nu q\bar{q}$ analyses, but additional modelling ones may be considered. For brevity, only areas where analyses differ in their treatment of systematics uncertainties with respect to the $\ell\nu q\bar{q}$ analysis will be discussed below. The list of systematics considered in each analyses can be seen explicitly in Tables B.1 - B.6.

Source	Corr	$\nu\nu qq$	$\ell\nu qq$	$\ell\ell qq$	$\ell\ell\nu\nu$	$\ell\nu\ell\nu$	$\ell\nu\ell\ell$	$qqbb$	$\nu\nu bb$	$\ell\nu bb$	$\ell\nu$
Small- R jet energy scale	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B
Small- R jet energy resolution	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B
Small- R jet flavor	Yes	S+B	S+B	S+B	S+B	S+B	–	S+B	S+B	S+B	S+B
Small- R jet pileup	Yes	S+B	S+B	S+B	S+B	S+B	–	S+B	S+B	S+B	S+B
Small- R jet punch-through	Yes	S+B	S+B	S+B	S+B	S+B	–	S+B	S+B	S+B	S+B
Small- R jet JVT	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B
Small- R jet Rtrk	Yes	–	–	–	–	–	–	S+B	S+B	S+B	–

Table B.1: Small- R jet systematic uncertainties. The abbreviations “S” and “B” stand for signal and background, respectively, while “–” denotes uncertainties that are either not applicable or are negligible, and “Corr” marks whether the uncertainty is correlated between the channels listed.

Source	Corr	$qqqq$	$\nu\nu qq$	$\ell\nu qq$	$\ell\ell qq$	$qqbb$	$\nu\nu bb$	$\ell\nu bb$	$\ell\ell bb$
Large- R jet D2R	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B
Large- R jet resolution	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B
Large- R jet JMR	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B
Large- R jet kinematics	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B
Large- R jet Rtrk	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B

Table B.2: Large- R jet systematic uncertainties. Same conventions as Table B.1 are used.

B.3.1 Experimental Uncertainties

Small- R Jets

Most analyses use a 21 parameter scheme for the JES while $\ell\ell\ell\ell$, $\ell\ell\nu\nu$, and $\ell\nu\ell\ell$ use the 3. All analyses use a 1-parameter scheme for the JER.

b -tagged efficiency

Most analysis implement b -tagged using small- R jets. The Vh analysis additionally use track jets to identify b -hadrons within a merged $h \rightarrow bb$ large- R jet. The same reduction scheme implementing three, four, five eigen-parameters for the b/c /light jets is used for the track-jet b -tagged uncertainties which are considered uncorrelated with the small- R ones. Tests were done where these collections of uncertainties were 100% correlated, which showed no visible impact in the fits or the results.

Source	Corr	$\nu\nu qq$	$\ell\nu qq$	$\ell\ell qq$	$\ell\ell\nu\nu$	$\ell\nu\ell\nu$	$\ell\nu\ell\ell$	$qqbb$	$\nu\nu bb$	$\ell\nu bb$	$\ell\ell bb$
b -tagging	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B
c -tagging	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B
light- q tagging	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B
tagging extrapolation	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B

Table B.3: Flavor-tagging systematic uncertainties. Same conventions as Table B.1 are used.

Source	Corr	$\ell\nu q\bar{q}$	$\ell\ell q\bar{q}$	$\ell\ell\nu\nu$	$\ell\nu\ell\nu$	$\ell\nu\ell\ell$	$\ell\ell\ell\ell$	$\ell\nu b\bar{b}$	$\ell\ell b\bar{b}$	$\ell\nu$	$\ell\ell$
Electron trigger	Yes	S+B	S+B	S+B	S+B	–	S+B	S+B	S+B	–	–
Electron reconstruction	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	–	–
Electron identification	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	–	S+B
Electron isolation	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	–	S+B
Electron energy scale	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B
Electron energy resolution	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	–	S+B
Electron fake background	No	–	–	–	–	–	–	–	–	B	B
Muon trigger	Yes	S+B	S+B	S+B	S+B	–	S+B	S+B	S+B	S+B	–
Muon reconstruction	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B
Muon isolation	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	–	S+B
Muon energy scale	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	–	–
Muon energy resolution	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B	S+B
Muon fake background	No	–	–	–	–	–	–	–	–	B	B

Table B.4: Lepton systematic uncertainties. Same conventions as Table B.1 are used.

Source	Corr	$\nu\nu q\bar{q}$	$\ell\nu q\bar{q}$	$\ell\ell\nu\nu$	$\ell\nu\ell\nu$	$\ell\nu\ell\ell$	$\ell\nu b\bar{b}$	$\ell\nu$
E_T^{miss} trigger	Yes	S+B	S+B	S+B	S+B	–	S+B	–
E_T^{miss} SoftTrk scale	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B
E_T^{miss} SoftTrk resolution	Yes	S+B	S+B	S+B	S+B	S+B	S+B	S+B

Table B.5: E_T^{miss} systematic uncertainties. Same conventions as Table B.1 are used.

Electrons and Muons

In addition to the nominal choice of lepton modeling uncertainties, the $\ell\nu\ell\nu$ channel considers an additional uncertainty on the electron identification.

B.3.2 Modeling Uncertainties

Majority of the differences in how uncertainties are handled are in the context of modeling uncertainties. This is due to several reasons, mainly due to the fact that different analysis have different background and analysis procedure, but also from the fact that there is no consistent way to treat all uncertainties of this type.

$t\bar{t}$ modelling

The VV and VH semileptonic channels evaluate the modeling and normalization uncertainty by varying several generator options and also comparing generators. For all of these analyses the default $t\bar{t}$ generator was POWHEG for showering and Pythia for hadronization. The following MC comparison were made and the differences taken as an uncertainty

- Two point matrix-element uncertainty by switching POWHEG with Madgraph
- Two point hadronization uncertainty by switching Pythia with HERWIG++
- Comparison with ISR/FSR enriched/depleted samples where the renormalization and factorization scales and the Powheg h_{damp} parameter is adjusted alongside a dedicated

Source	Corr	$\nu\nu qq$	$\ell\nu qq$	$\ell\ell qq$	$\ell\ell\nu\nu$	$\ell\nu\ell\nu$	$\ell\nu\ell\ell$	$qqbb$	$\nu\nu bb$	$\ell\nu bb$	$\ell\ell bb$	$\ell\nu$	$\ell\ell$
DY PDF variation	Yes	–	–	–	–	–	–	–	–	–	–	B	B
DY PDF choice	Yes	–	–	–	–	–	–	–	–	–	–	B	B
DY PDF scale	Yes	–	–	–	–	–	–	–	–	–	–	negl.	B
DY α_s	Yes	–	–	–	–	–	–	–	–	–	–	B	B
DY EW corrections	Yes	–	–	–	–	–	–	–	–	–	–	B	B
DY photon-induced	Yes	–	–	–	–	–	–	–	–	–	–	–	B
Top cross section	No	B	F	F	B	B	–	B	B	B	B	B	negl.
Top extrapolation	No	–	–	–	–	–	–	–	–	–	–	B	–
Top modeling	No	B	B	B	B	B	–	–	B	B	B	negl.	negl.
Diboson cross section	No	B	B	B	B	B	–	B	B	B	B	negl.	negl.
Diboson extrapolation	No	–	–	–	–	–	–	–	–	–	–	B	–
Multijet cross section	No	–	B	–	–	–	B	B	–	B	–	–	–
Multijet modeling	No	–	–	–	–	–	–	B	–	B	–	B	B
Z+jets cross section	No	F	B	F	–	–	–	–	B	B	B	–	–
Z+jets modeling	No	B	B	B	–	–	–	–	B	B	B	–	–
W+jets cross section	No	B	F	B	–	–	–	–	B	B	B	–	–
W+jets modeling	No	B	B	B	–	–	–	–	B	B	B	–	–

Table B.6: Theoretical systematic uncertainties. Same conventions as Table B.1 are used with the addition that the abbreviation “F” means that this parameter was left to float in the background control region for that channel.

change tune.

Other analyses with non-negligible $t\bar{t}$ backgrounds only considered normalization uncertainties which was measured to be approximetly 30% [258].

V+jets modelling

In the $\nu\nu qq$, $\ell\nu qq$ and the semileptonic VH channels the uncertainty on the V+jets modeling is obtained by comparing the nominal shape of the discriminant variable obtained with SHERPA to the one obtained with an alternative samples. The list of variations considered as uncertainties are:

- Two-point generator comparison with respect to Madgraph +Pythia
- 7 Scale variations of the renormalization and factorization scale
- 2 CKKW matching scale and 2 Resummation scale variations
 - Only evaluated in $\ell\nu qq$
- 2 NNPDF3.0nnlo α_s variations
- 100 NNPDF3.0nnlo variations [121]

The $\ell\ell qq$ uses a data-driven approach to estimate the uncertainties known as the α -ratio method. This is done by comparing the MC in the signal region by extrapolated data from a high-purity control region. The data in the control region is extrapolated by a term $\alpha(x)$ which is the ratio of sample MC in the signal region to the control region.

Diboson modelling

The semileptonic VV and VH channels consider an overall 30% normalization uncertainty which conservatively includes the $\sim 10\%$ expected from cross-section uncertainty and $\sim 20\%$ expected from normalization and factorization scale uncertainties.

The leptonic $\ell\ell\ell$, and $\ell\ell\nu\nu$ consider additional uncertainties on the diboson modeling associated to a functional fit to the background. These channels also consider additional theoretical normalization uncertainties due to both $gg \rightarrow ZZ$ and $q\bar{q} \rightarrow ZZ$ contributions.

Multijet modeling

The multijet contamination is evaluated differently in several analyses. The $\ell\nu qq$ channel evaluated the multijet contamination with a fake-factor method. Systematic uncertainties are evaluated for possible variations in the electroweak background, fake-flavour composition, energy scale, and isolation. The $\ell\nu bb$ channel instead uses a template-method for the fake lepton isolation and the normalization and shape uncertainty are directly estimated for the evaluated template.

For the $q\bar{q}bb$ analyses the normalization and shape of the multi-jet contribution are evaluated in validation regions by varying empirical fits to the data. In the $qqqq$ channel, multijet is the only relevant background and estimated directly from the final fit.

Signal Uncertainty

The main uncertainties in regards to the signal samples is the effect of PDF uncertainties and possible ISR/FSR variations. PDF variations can effect the normalization of signals and ISR/FSR can effect the signal acceptance due to different jet multiplicity for example.

The PDF uncertainties for VV channels are estimated by taking the acceptance difference between the NNPDF23LO [122] and MSTW2008LO [214] PDFs and adding it in quadrature to the difference between the NNPDF23LO error sets. A conservative estimate of 2% normalization uncertainty is applied. Similarly the difference between varied Pythia [200] tunes for ISR/FSR are used to derive normalization uncertainties of $< 3\%$.

For the dilepton channels the PDF variation uncertainty is obtained using the 90% CL CT14NNLO PDF [124] error set. Rather than using a single nuisance parameter to describe the 28 eigenvectors of this PDF error set, which could lead to an underestimation, a re-diagonalised set of 7 PDF eigenvectors was used, and treated as correlated.

Appendix C

Glossary

ALICE	A Large Ion Collider Experiment
ATLAS	A Toroidal LHC Apparatus
aQGC	Anomalous Quartic Gauge Couplings
BCM	Beam Conditions Monitor
BDT	Boosted Decision Tree
BSM	Beyond the Standard Model
CERN	The European Organization for Nuclear Research
CMS	Compact Muon Solenoid Experiment
CSC	Cathode Strip Chamber
DQ	Data Quality
EFT	Effective Field Theory
EM	Electromagnetic
EMEC	Electromagnetic End-cap Calorimeter
EMTopo	Electromagnetically Calibrated Topologically Clustered
FCal	Forward Calorimeter
ggF	Gluon-Gluon Fusion
HCal	Hadronic Calorimeter
HEC	Hadronic End-cap Calorimeter
HVT	Heavy Vector Triplet
HLT	High Level Trigger
HL-LHC	High-Luminosity Large Hadron Collider
ID	Inner Detector
IBL	Insertable B-Layer
IP	Interaction Point
JES	Jet Energy Scale
JER	Jet Energy Resolution
JMS	Jet Mass Scale
JMR	Jet Mass Resolution

JVT	Jet Vertex Tagger
LHC	Large Hadron Collider
LHCb	Large Hadron Collider Beauty Experiment
L1	Level-One Trigger
LAr	Liquid-Argon
LCW	Local Cluster Re-weighting
LCTopo	Locally Calibrated Topologically Clustered
MDT	Monitored Drift Tubes
MC	Monte-Carlo
MS	Muon Spectrometer
PDF	Parton Distribution Function
PV	Primary Vertex
pdf	Probability Distribution Function
QCD	Quantum Chromo-Dynamic
QFT	Quantum Field Theory
$q\bar{q}$	Quark-antiquark Fusion
RS	Randall-Sundrum
RPC	Resistive plate chambers
SCT	Silicon Micro-strip Tracker
SM	Standard Model
TGC	Thin Gap Chambers
Tile	Tile Calorimeter
TCC	Track-Calo Cluster
TRT	Transition Radiation Tracker
TDAQ	Trigger and Data Acquisition System
VBF	Vector-boson Fusion

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