A Search for a New Particle Decaying to Pairs of Weak Gauge Bosons with the ATLAS Detector

by

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Dissertation submitted in partial fulfillment of the requirements for the degree of Doctor of Philosophy in the Department of Physics in the Graduate School of Duke University
2016
ABSTRACT

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Abstract

A search for new heavy resonances decaying to boson pairs (WZ, WW or ZZ) using 20.3 fb$^{-1}$ of proton-proton collision data at a center of mass energy of 8 TeV is presented. The data were recorded by the ATLAS detector at the Large Hadron Collider (LHC) in 2012. The analysis combines several search channels with the leptonic, semi-leptonic and fully hadronic final states. The diboson invariant mass spectrum is studied for local excesses above the Standard Model background prediction, and no significant excess is observed for the combined analysis. 95% confidence limits are set on the cross section times branching ratios for three signal models: an extended gauge model with a heavy W boson, a bulk Randall-Sundrum model with a spin-2 graviton, and a simplified model with a heavy vector triplet. Among the individual search channels, the fully-hadronic channel is predominantly presented where boson tagging technique and jet substructure cuts are used. Local excesses are found in the dijet mass distribution around 2 TeV, leading to a global significance of 2.5 standard deviations. This deviation from the Standard Model prediction results in many theory explanations, and the possibilities could be further explored using the LHC Run 2 data.
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1
Introduction

The Standard Model has been a huge success in describing fundamental particles in observable matter and their interactions. The newly-discovered Higgs boson adds more credit to how well the Standard Model predicts the nature of the universe.

Despite its success, the Standard Model is imperfect in many ways such as the lack of explanation for the dark matter/dark energy, that it does not include the gravitational force, or solve the hierarchy problem. Therefore, there exist many other theories beyond Standard Model and searches for new physics which could possibly solve those questions not answered by the Standard Model. One way of searching for new physics is to look for potential heavy particles with certain decay modes and to look at the reconstructed invariant mass spectrum of its decay products. If there were a new particle, a bump would appear in the mass spectrum. This thesis presents a search for such a narrow resonance decaying to boson pairs ($WZ$, $WW$ and $ZZ$). The search with fully hadronic final states is focused on since they provide a much larger branching fraction than the leptonic and semileptonic decay modes and also make it possible to extend the search region to the highest possible mass range given current collision data. The result of combining all decay modes are also
shown, providing a better sensitivity within a wide mass range.

The contents of this document are organized as follows: the first few chapters introduce the necessary background information and the last two chapters talk about the analyses. Chapter 2 introduces the theoretical framework. The experimental setup is discussed in Chapter 3: the LHC and the ATLAS detector. The Monte Carlo technique used to simulate the proton-proton collisions is described in Chapter 4, followed by the discussion of event and object reconstruction in Chapter 5. The focus of this thesis is presented in Chapter 6 and 7, in which a search for new heavy resonances decaying to boson pairs is presented.
2.1 Introduction to the Standard Model

2.1.1 Overview

Our understanding of the fundamental particles of the universe and their interactions is currently described by a theoretical framework called Standard Model (SM). Mathematically speaking, it is a quantum field theory (QFT) based on an $SU(3) \otimes SU_L(2) \otimes U(1)$ gauge symmetry. Figure 2.1 introduces the structure of the SM: it contains fermions (with half integer spin which follow the Pauli exclusion principle) and bosons (with integer spin which follow Bose-Einstein statistics). In other words, the SM has the structure of building blocks of matter (fermions) and the mediators of forces which control their interactions (bosons).

Fermions are further divided into quarks, which have color and interact via strong force, and leptons, which do not. Both quarks and leptons have three generations, grouped generally based on their increasing masses. Quarks are further divided into two categories: up-type and down-type, after the name of the first generation. Leptons also exist in two categories: charged leptons and neutrinos. The masses of the
Table 2.1: Summary of the gauge bosons mediating fundamental interactions.

<table>
<thead>
<tr>
<th>Gauge Boson</th>
<th>Charge ($e$)</th>
<th>Spin</th>
<th>Force mediated</th>
<th>Mass (GeV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\gamma$</td>
<td>0</td>
<td>1</td>
<td>EM</td>
<td>0</td>
</tr>
<tr>
<td>$W^\pm$</td>
<td>$\pm 1$</td>
<td>1</td>
<td>Weak</td>
<td>80.398±0.025</td>
</tr>
<tr>
<td>$Z^0$</td>
<td>0</td>
<td>1</td>
<td>Weak</td>
<td>91.1876±0.0021</td>
</tr>
<tr>
<td>gluons</td>
<td>0</td>
<td>1</td>
<td>Strong</td>
<td>0</td>
</tr>
<tr>
<td>graviton (hypothetical)</td>
<td>0</td>
<td>2</td>
<td>Gravity</td>
<td>0</td>
</tr>
</tbody>
</table>

neutrinos are assumed to be zero by the SM, however, their masses are observed to be nonzero though very small in reality. Each of these particles have their antiparticles with the same mass but opposite charge.

Figure 2.1: The Standard Model [1]

Bosons mediate the electromagnetic, weak and strong forces described in the SM. Photons mediate the electromagnetic interactions, and follow the description of the
Table 2.2: Summary of the fermions

<table>
<thead>
<tr>
<th>Fermions</th>
<th>charge</th>
<th>spin</th>
<th>Mass</th>
</tr>
</thead>
<tbody>
<tr>
<td>u</td>
<td>$+\frac{2}{3}$</td>
<td>$\frac{1}{2}$</td>
<td>$1.5 \sim 3.3$ MeV</td>
</tr>
<tr>
<td>d</td>
<td>$-\frac{1}{3}$</td>
<td>$\frac{1}{2}$</td>
<td>$3.5 \sim 6.0$ MeV</td>
</tr>
<tr>
<td>c</td>
<td>$+\frac{2}{3}$</td>
<td>$\frac{1}{2}$</td>
<td>$1.27^{+0.07}_{-0.11}$ GeV</td>
</tr>
<tr>
<td>s</td>
<td>$-\frac{1}{3}$</td>
<td>$\frac{1}{2}$</td>
<td>$104^{+26}_{-34}$ MeV</td>
</tr>
<tr>
<td>t</td>
<td>$+\frac{2}{3}$</td>
<td>$\frac{1}{2}$</td>
<td>$171.2 \pm 2.1$ GeV</td>
</tr>
<tr>
<td>b</td>
<td>$-\frac{1}{3}$</td>
<td>$\frac{1}{2}$</td>
<td>$4.20^{+0.17}_{-0.07}$ GeV</td>
</tr>
<tr>
<td>e</td>
<td>-1</td>
<td>$\frac{1}{2}$</td>
<td>$0.510998910 \pm 1.3e-8$ MeV</td>
</tr>
<tr>
<td>$\nu_e$</td>
<td>0</td>
<td>$\frac{1}{2}$</td>
<td>$&lt; 2$ eV</td>
</tr>
<tr>
<td>$\mu$</td>
<td>-1</td>
<td>$\frac{1}{2}$</td>
<td>$105.658367 \pm 4e-6$ MeV</td>
</tr>
<tr>
<td>$\nu_\mu$</td>
<td>0</td>
<td>$\frac{1}{2}$</td>
<td>$&lt; 0.19$ MeV</td>
</tr>
<tr>
<td>$\tau$</td>
<td>-1</td>
<td>$\frac{1}{2}$</td>
<td>$1776.84 \pm 0.17$ MeV</td>
</tr>
<tr>
<td>$\nu_\tau$</td>
<td>0</td>
<td>$\frac{1}{2}$</td>
<td>$&lt; 18.2$ MeV</td>
</tr>
</tbody>
</table>

theory of quantum electrodynamics (QED). The weak force is carried by $W^+, W^-$ and $Z^0$ bosons. Unlike massless photons, these are massive particles. The reason for this difference will be explained in section 2.1.4. Gluons mediate the strong forces between color charged particles (quarks and themselves). Gluons are massless, and their interactions are described by the theory of quantum chromodynamics (QCD).

2.1.2 Quantum Chromodynamics

The strong interaction is described by QCD with a $SU(3)_c$ gauge symmetry on quark color fields. Quarks exist as color triplets, carrying colors red, blue or green; gluons exist as color octet, transferring quarks to different color states during quark-gluon interactions with coupling strength $g_S$. Leptons and electroweak bosons are colorless particles which do not interact with gluons.
The $SU(3)_c$ gauge invariant Lagrangian can be written as,

$$\mathcal{L}_{QCD} = \bar{q}_f (i \gamma^\mu D_\mu - m_f) q_f - \frac{1}{4} G^a_{\mu\nu} G^a_{\mu\nu}$$

(2.1)

where the covariant derivative and the gauge field strength are,

$$D_\mu \equiv \partial_\mu - ig_s \frac{\lambda_a}{2} G^a_\mu$$

$$G^a_{\mu\nu} \equiv \partial_\mu G^a_\nu - \partial_\nu G^a_\mu + g_s f^{abc} G^b_\mu G^c_\nu$$

(2.2)

$G^a_\mu$ are the gluon fields and the coupling strength is $g_s$. The fields $G$ and $q$ are functions of spacetime. $f^{abc}$ are the structure constants, $abc$ are color indices, and $\mu\nu$ are spacetime indices. If the Lagrangian $\mathcal{L}_{QCD}$ is expanded, one gets the color interaction between the quark fields and the gluon fields as well as the cubic and quartic gluon self-interaction due to the non-abelian character of the color group.

Based on the knowledge that no free colored particles are found in the natural world, one has to make the assumption that all hadrons have to be colorless. The color singlet for mesons ($q\bar{q}$) is $\Psi = \frac{1}{\sqrt{3}} (r\bar{r} + b\bar{b} + g\bar{g})$ and for baryons, $\Psi = \frac{1}{\sqrt{6}} (rgb - rbg + grb - bgr)$, where $r =$ red, $b =$ blue and $g =$ green.

It is hypothesized that the quarks have to be confined within color singlet bound states, which is known as the confinement hypothesis. Nevertheless, the evolution of the QCD coupling $\alpha_s$ which defines the strength of an interaction between two color charges, described below, implies that the confined quarks are “asymptotically free”, which means that the quarks interact weakly at high energies.

In spite of the simple $\mathcal{L}_{QCD}$ induced by the color symmetry, the physics in it is quite rich. Since all the interactions in $\mathcal{L}_{QCD}$ are given by one single universal coupling $g_s$, $\alpha_s$ should be expressed as a function of $g_s$ (usually written as $\alpha_s = \frac{g_s}{4\pi}$). Actually, $\alpha_s$ is not a constant, but depends on the energy scale of the interaction:

$$\alpha_s(Q^2) = \frac{2\pi}{-\beta_0 \ln (\frac{\mu^2}{\Lambda^2})}$$

(2.3)
Figure 2.2: Measurements of $\alpha_s$ as a function of the energy scale $Q$ [3].

where $\beta_0 = 1/3 \times N_f - 11/2$ [2]. $\alpha_s$ is close to 1 at low energies, so perturbation theory is no longer valid. It goes to zero if $\mu$ is much larger than the energy scale $\Lambda$, indicating that perturbation theory is recovered at high energy. Figure 2.2 shows the experimental evidence of $\alpha_s$ depending on the scale energy.

2.1.3 Electroweak Interactions

The Standard Model electroweak theory is governed by the $SU(2)_L \otimes U(1)$ gauge symmetry that characterizes the interactions of the electroweak vector bosons with quarks and leptons.

The electroweak Lagrangian can be expressed,

$$\mathcal{L}_{EW} = \bar{\psi} i\gamma^\mu D_\mu \psi - \frac{1}{4} F_{\mu\nu}^i F_{\mu\nu}^i - \frac{1}{4} W_{\mu\nu} W^{\mu\nu} \quad (2.4)$$

where $\psi$ indicates the spinors of fermions including the quarks, the leptons and the neutrinos.

$D_\mu$ is the covariant derivative,

$$D_\mu \equiv \partial_\mu + i\frac{g}{2} T_i A_\mu^i + i\frac{g'}{2} Y B_\mu \quad (2.5)$$

where $g$ and $g'$ are the gauge couplings.
Table 2.3: Strength of the Z boson’s vector (V) and axial vector (A) couplings to fermions

<table>
<thead>
<tr>
<th>fermion</th>
<th>$t_3$</th>
<th>$Q$</th>
<th>$g_v$</th>
<th>$g_a$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\nu_e, \nu_\mu, \nu_\tau$</td>
<td>$\frac{1}{2}$</td>
<td>0</td>
<td>$\frac{1}{2}$</td>
<td>$\frac{1}{2}$</td>
</tr>
<tr>
<td>$e, \mu, \tau$</td>
<td>$\frac{1}{2}$</td>
<td>$-1$</td>
<td>$\frac{1}{2} + 2\sin^2\theta_w$</td>
<td>$\frac{1}{2}$</td>
</tr>
<tr>
<td>$u, c, t$</td>
<td>$\frac{1}{2}$</td>
<td>$\frac{1}{2}$</td>
<td>$\frac{1}{2} - \frac{1}{2}\sin^2\theta_w$</td>
<td>$\frac{1}{2}$</td>
</tr>
<tr>
<td>$d, s, b$</td>
<td>$\frac{1}{2}$</td>
<td>$-\frac{1}{2}$</td>
<td>$\frac{1}{2} + \frac{1}{2}\sin^2\theta_w$</td>
<td>$\frac{1}{2}$</td>
</tr>
</tbody>
</table>

The field tensors with the gauge fields ($A^i_\mu$ with $i=1,2,3$ and $B_\mu$) corresponding to the group generators are,

\[
F^i_{\mu\nu} \equiv \partial_\mu A^i_\nu - \partial_\nu A^i_\mu + g\epsilon^{ijk}A^j_\mu A^k_\nu
\]
\[
G_{\mu\nu} \equiv \partial_\mu B_\nu - \partial_\nu B_\mu
\]  

(2.6)

Gauge symmetry does not allow an explicit mass term either for fermions or for gauge bosons in the electroweak Lagrangian. The $L_{EW}$ can be separated into three currents. There is a charged current term that is mediated by the $W$ bosons, which couples only to left-handed fermions due to the vector-axial (V-A) structure of the coupling. Electromagnetic and weak neutral current are mediated by the photons and $Z$ bosons respectively. The electromagnetic interaction couples to all particles with electric charge, while weak neutral current interactions have some mixing of vector and axial couplings. The strength of the $Z$ boson’s V-A couplings to fermions, at tree level in Standard Model, are given in Table 2.3.

2.1.4 The Higgs Mechanism

The previous formalism predicts the existence of four massless mediators to describe the electroweak force. However, the $W$ and $Z$ bosons, associated with the weak force are massive. Moreover, fermions must be described as massless particles, contrary to observation. The mass terms for fermions and bosons break the invariance of the Lagrange density under gauge transformation. To solve this problem, the Higgs
mechanism is introduced, which describes the spontaneous breaking of the symmetries to generate mass for $W^\pm$ bosons and the $Z^0$ boson, while keeping the photon massless [4] [5] [6].

The Higgs scalar field is introduced in order to apply the Higgs mechanism to electroweak theory,

$$L_{\text{scalar}} = \partial_\mu \Phi^\dagger \partial^\mu \Phi - V(\Phi^\dagger \Phi)$$

(2.7)

where a complex doublet is introduced in this simplest case,

$$\Phi \equiv (\phi^+ \phi)$$

(2.8)

with a particular potential,

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

(2.9)

where $\mu$ and $\lambda$ are real constants.

Different choices of $\mu$ and $\lambda$ can be made when searching for the minimum of the potential. The case $\lambda < 0$ is unphysical, leading to unstable minimum. When $\lambda > 0$, two possible solutions are given. If $\mu^2 > 0$, a vacuum state $\Phi = 0$ arises, where the symmetry stays unbroken. If $\mu^2 < 0$, the symmetry breaks down spontaneously and the original symmetry group $SU(2)_L \otimes U(1)_Y$ is broken to $U(1)_{EM}$. Figure 2.3 shows the potential as a function of the field for the two possible signs of $\mu$. The vacuum expectation of the Higgs field can, without loss of generality, be chosen as,

$$<\Phi>_0 = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v \end{pmatrix}$$

(2.10)

where $v = \sqrt{-\frac{\mu^2}{\lambda}}$ is the vacuum expectation value.

The terms generated by symmetry breaking that describe the new boson masses are

$$L_{\text{mass}} = \frac{g^2 v^2}{4} W_\mu^+ W^{-\mu} + \frac{g^2 v^2}{8 \cos^2 \theta_w} Z_\mu Z^\mu - \frac{1}{2} (2 v^2 \lambda) H^2$$

(2.11)
With the $\mathcal{L}_{mass}$, the masses associated with the electroweak mediators as well as the Higgs boson mass can be verified,

$$M_W = \frac{g v}{2}, \quad M_Z = \frac{g v}{2 c_w} = \frac{M_W}{c_w}, \quad M_H = \sqrt{2 v^2 \lambda}$$

(2.12)

![Figure 2.3: The potential $V$ of the scalar field in the case $\lambda > 0$ and $\mu^2 > 0$ (left) or $\mu^2 < 0$ (right) [7].](image)

To summarize, the electromagnetic and the weak forces are unified as an “electroweak” force. The vector bosons get their masses via the Higgs mechanism which is also known as the electroweak symmetry breaking. The Higgs boson, which carries the Higgs field in the Higgs mechanism, is the final piece of the puzzle of the SM. In 2012, both ATLAS and CMS experiments claimed that they have observed a new particle that is consistent with the Higgs boson predicted by the SM [8]. Although this discovery puts the mechanism for electroweak symmetry breaking in nature on solid ground, it leads to the further question of why the Higgs mass (measured to be around 125 GeV by ATLAS and CMS) is so small despite the large loop corrections to this parameter. The question is referred to as the hierarchy problem.
2.2 Physics of Hadron Collisions

The SM processes at hadron colliders most important to understanding the measurements presented in this thesis are briefly presented here.

The phenomena observed in hadron collisions are dominated by QCD. One of the main features of it is the running of the coupling constant $\alpha_S$. As described in Section 2.1.2, at high momentum transfer, $\alpha_S$ gets very small and the theory may be used in the perturbative regime where precise calculations can be performed; at low momentum transfer, $\alpha_S$ gets bigger and perturbative calculations fail. In hadron collisions, a very small fraction of events can be described in perturbative regime while most of the events are at low momentum transfer.

2.2.1 Low-$Q^2$ QCD Events

The vast majority of the hadron collisions at LHC are not parton scattering with large momentum transfer (although LHC was designed to probe physics at center-of-mass energies above 10 GeV). The effect of the soft interactions, therefore, must be taken into account to describe the reconstruction of the high-energy collisions well. These soft QCD processes include interactions between the remnant partons (called the underlying event). In addition, since the most likely proton-proton scattering process is soft, multiple hadron-hadron interactions due to the high luminosity of the LHC (called pileup) are dominated by non-perturbative QCD processes. The scale that sets apart the soft and hard interactions is $\Lambda_{QCD}$ (the scale defined in section 2.1.2), which is about 200 MeV.

2.2.2 Hard QCD Interactions

At high momentum transfer, the interactions of protons can be treated as direct interactions of the internal partons, and the protons’ internal structures need to be understood to correctly model the process. This understanding is mainly achieved by
Deep Inelastic Scattering (DIS) experiments [9] [10] [11] in which a lepton interacts directly with a parton inside the proton. The interaction cross section and scattered lepton kinematics as a function of the momentum transferred are sensitive to the proton’s structure. Parton distribution functions (PDFs) are used to describe the probability of a parton of a certain flavor to carry the fractional momentum $x$ of the proton at a momentum transfer $Q^2$. PDFs can be measured in Deep Inelastic Scattering (DIS) experiments and their evolution as described by theory (and checked by experiment) is used to predict physical cross sections at the LHC. Although the PDFs themselves cannot be calculated analytically and have to be measured, their evolution in $Q^2$ can be determined using splitting functions, formally written down in the DGLAP-equations [12] [13]. With the DGLAP equations to predict the evolution of structure functions combined with experimental measurements at certain data points, a full picture of the PDFs can be formed.

2.2.3 Diboson Production and Decay Modes

In the SM, diboson production proceeds principally through the quark anti-quark annihilation process and gluon-gluon fusion process at leading order. The SM measurements of the diboson cross section are roughly 100 pb for $WW$, 20 pb for $WZ$ and 10 pb for $ZZ$, which are very small compared to $10^4$ pb/GeV for dijet differential cross section at dijet mass $\sim 100$ GeV.

$W$ and $Z$ bosons decay either hadronically or leptonically, and their branching fractions are showed in Table 2.4. Approximately 70% of $W$ and $Z$ decays result in a quark and antiquark. The other decays are to charged leptons and (or) neutrinos.
Table 2.4: $W$ and $Z$ boson branching fractions

<table>
<thead>
<tr>
<th>$W$ Decay Mode</th>
<th>$W$ Decay Fraction</th>
<th>$Z$ Decay Mode</th>
<th>$Z$ Decay Fraction</th>
</tr>
</thead>
<tbody>
<tr>
<td>$qq'$</td>
<td>67%</td>
<td>$qq$</td>
<td>70%</td>
</tr>
<tr>
<td>$l\nu$</td>
<td>22%</td>
<td>$ll$</td>
<td>7%</td>
</tr>
<tr>
<td>$\tau\nu$</td>
<td>11%</td>
<td>$\tau\tau$</td>
<td>3%</td>
</tr>
<tr>
<td>$\nu\nu$</td>
<td>-</td>
<td>$\nu\nu$</td>
<td>20%</td>
</tr>
</tbody>
</table>

2.3 Physics of Jets

2.3.1 Hadronization and Jet

After high-energy collisions in colliders, free quarks or gluons are created. The accelerated colored partons will emit QCD radiation in the form of gluons, resulting in a shower of partons. Due to color confinement, these cannot propagate long distances; after the parton shower, remaining quarks are combined spontaneously to form hadrons. We call this process hadronization.

Once the hadrons are detected, how we measure them is very important. Usually they are combined into jets, which ideally contain all the hadrons arising from one parton. We then use the jet to investigate the hard process that produced it. This can be done by a well defined jet algorithm.

2.3.2 Jet Algorithms

A jet algorithm is a method to combine objects into jets which mimic the partons that produce them. When we come up with a jet algorithm, one of the most important factors is how to decide which objects should be included in a jet and how to combine their momenta meaningfully.

The simplest jet algorithms are the cone algorithms which start from one particle, and take all the constituents around it which are separated in $\sqrt{(\Delta \eta)^2 + (\Delta \phi)^2}$ by less than the jet radius $R$. This is done iteratively by choosing a seed particle,
combining the momenta of all the particles nearby within the cone and redefining the axis. At first, the axis almost surely changes at each step. The iterative step does not stop until the axis is stable. Clearly, the jet obtained this way is independent of the choice of the seed. The biggest disadvantage of this algorithm is that it is easily affected by soft emissions, for instance, it falls into trouble when two nearby cones are overlapped since it has to decide whether the two cones should be merged and which soft particles go to which cone. Therefore, this kind of jet algorithm is not infrared safe with the exception of SIScone [14] which takes into account all the possible scenarios.

The most commonly used method of forming jets nowadays is successive recombination algorithms. The basic idea is to attempt to go backwards through the parton shower by merging two particles at each step. The closest particles are chosen to be merged at each stage, with their distance defined as:

\[ d_{ij}^2 = \min(p_{T_i}^{2\beta}, p_{T_j}^{2\beta}) \frac{\Delta R_{ij}^2}{R^2} \]  

(2.13)

\[ d_{iB} = p_{T_i}^{2\beta} \]  

(2.14)

where \( d_{ij} \) is the distance between particle \( i \) and \( j \), \( d_{iB} \) is introduced to control when merging stops, and is historically referred to the distance between particle \( i \) and the beamline, \( R \) is the jet cone size chosen before running the algorithm and \( \Delta R_{ij}^2 = (\eta_i^2 - \eta_j^2)^2 + (\phi_i^2 - \phi_j^2)^2 \). The distances \( d_{ij} \) and \( d_{iB} \) are calculated for all particles, then pair of particles corresponding to the smallest value in distance are merged by summing their four-momenta. If the distance \( d_{iB} \) is the smallest, then \( i \) is defined as a jet and removed from the clustering procedure. The previous steps are repeated until all the particles are merged into jets.

In the above equation, \( \beta \) can be assigned different values resulting in different
algorithms. If $\beta$ is assigned zero, we call the algorithm the Cambridge-Aachen (CA) algorithm [15] which obviously is dependent only on the angular information in the distance calculation. The jet shapes produced by this algorithm are irregular; however, when comes to substructure, CA jets can be useful for characterizing structure at a variety of angular scales. If $\beta$ is assigned -1, we have anti-$k_T$ algorithm [16], which is the most widely used algorithm in ATLAS. It starts from the particles with the highest $p_T$ leading to a hard center surrounded with soft constituents. The jets reconstructed this way turn out to be circular as a result. If $\beta$ is assigned 1, we have $k_T$ algorithm [17] which clusters soft particles first resulting in irregular shaped jets.

Figure 2.4 shows the jet displays for the same event in the $\eta - \phi$ space for the Cambridge-Aachen, anti-$k_T$, $k_T$ and SIScone algorithms. As can be seen, all of the algorithms agree with the position of the hardest jets, although slight differences
exist in terms of which soft constituents to include. The circular characteristic shape of the anti-\(k_T\) jets can be observed from the figure.

2.4 Models with Heavy Resonances Decaying to Diboson

2.4.1 EGM \(W'\)

One of the main goals of colliders such as the LHC is to establish the gauge group which fully describes the strong and electroweak interactions. A series of precision measurements conducted up until now have probed physics at GeV and TeV scale, supporting the hypothesis that the group is of the form \(SU(3)_c \times SU(2)_L \times U(1)_Y\) as predicted by the SM. However, more complex gauge structures are proposed to exist at higher energies by many models. If any of these ideas were correct, LHC or other future colliders could find direct evidence for the existence of the new gauge structures. The extended gauge model (EGM) [18] provides a framework for describing a new heavy vector boson which serves as the partner of SM \(W\) and \(Z\) bosons.

The discovery of a new gauge boson, such as a \(W'\), would be a clear signature for new physics beyond the SM. In the EGM, this charged \(W'\) boson couples to SM fermions in a same way as \(W\) and \(Z\) bosons do, and \(W'WZ\) coupling has the same structure as \(WWZ\) coupling in the SM. The only difference is that \(W'WZ\) coupling is suppressed with respect to SM \(WWZ\) coupling by a factor of \(c(m_W/m_{W'})^2\) where \(m_W\) is the \(W\) boson mass, \(m_{W'}\) is the \(W'\) mass, and \(c\) is a scaling constant of order one which is taken to be unity here. The width of the \(W'\) mass becomes narrower due to this scaling factor.

Information on the cross sections, widths, and branching ratios for EGM \(W'\) production at \(\sqrt{s} = 8\) TeV are provided in Table 2.5. The cross section for EGM \(W'\) production is shown as a blue line in Figure 2.6. Leading order cross sections are used throughout the analysis in later chapters. Next-to-next-to-leading order K-factors are calculated using ZWProd [19], which are not used in this thesis.
Figure 2.5: Branching ratios for EGM $W'$

Figure 2.6: Cross section times branching ratio for production of a new heavy boson decaying to vector bosons. In the EGM $W'$ case, $c = 1$. For bulk $G^*$ production, $\bar{M}_{Pl} = 1$. 
Table 2.5: Cross section, width, and branching ratios for EGM $W'$ production at $\sqrt{s} = 8$ TeV.

<table>
<thead>
<tr>
<th>Pole Mass [TeV]</th>
<th>$\sigma \times B(W' \rightarrow WZ)$ [fb]</th>
<th>K-factor</th>
<th>Width [GeV]</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.20</td>
<td>49920.</td>
<td>1.305</td>
<td>5.5</td>
</tr>
<tr>
<td>0.25</td>
<td>42480.</td>
<td>1.310</td>
<td>7.5</td>
</tr>
<tr>
<td>0.30</td>
<td>24120.</td>
<td>1.311</td>
<td>9.5</td>
</tr>
<tr>
<td>0.35</td>
<td>13470.</td>
<td>1.310</td>
<td>11.0</td>
</tr>
<tr>
<td>0.40</td>
<td>7900.</td>
<td>1.307</td>
<td>13.0</td>
</tr>
<tr>
<td>0.60</td>
<td>1470.</td>
<td>1.289</td>
<td>21.0</td>
</tr>
<tr>
<td>0.80</td>
<td>415.0</td>
<td>1.269</td>
<td>28.0</td>
</tr>
<tr>
<td>1.00</td>
<td>149.6</td>
<td>1.251</td>
<td>35.0</td>
</tr>
<tr>
<td>1.20</td>
<td>61.30</td>
<td>1.235</td>
<td>43.0</td>
</tr>
<tr>
<td>1.40</td>
<td>27.00</td>
<td>1.220</td>
<td>50.0</td>
</tr>
<tr>
<td>1.60</td>
<td>12.61</td>
<td>1.206</td>
<td>58.0</td>
</tr>
<tr>
<td>1.80</td>
<td>6.156</td>
<td>1.192</td>
<td>65.0</td>
</tr>
<tr>
<td>2.00</td>
<td>3.109</td>
<td>1.176</td>
<td>72.0</td>
</tr>
<tr>
<td>2.20</td>
<td>1.646</td>
<td>1.161</td>
<td>80.0</td>
</tr>
<tr>
<td>2.40</td>
<td>0.8988</td>
<td>1.148</td>
<td>87.0</td>
</tr>
<tr>
<td>2.60</td>
<td>0.5066</td>
<td>1.139</td>
<td>94.0</td>
</tr>
<tr>
<td>2.80</td>
<td>0.2990</td>
<td>1.137</td>
<td>102.0</td>
</tr>
<tr>
<td>3.00</td>
<td>0.1861</td>
<td>1.143</td>
<td>109.0</td>
</tr>
</tbody>
</table>

2.4.2 Graviton

In the Standard Model theory, gauge fields exist in the traditional four dimensions of spacetime. However, the possibility of other compact spatial dimensions has not been ruled out by current experiments. Theodor Kaluza and Oskar Klein [20] first came up with the possible existence of extra dimensions of space in the middle of the 1920s as a means of unifying the electromagnetic and gravitational fields as components of a single higher dimensional field. In the theories of extra dimensions, the SM gauge fields are allowed to propagate into the bulk of the extra dimensions, manifesting themselves as apparent additional mass. In this case, the masses of the excited states will form highly degenerate towers, and the mass spectrum will be discrete if the dimension is finite, or compact. The discrete modes are defined as KK
...excitations.

Randall-Sundrum (RS) models [21] addresses the hierarchy problem by introducing warped extra dimensions compactified on a circle whose upper and lower halves are identified.

The background metric in the Randall-Sundrum model is parametrized by

\[ ds^2 = e^{-2k|y|} \eta_{\mu\nu} dx^\mu dx^\nu - dy^2 \]  \hspace{1cm} (2.15)

where \( k \) is the curvature of the warped extra dimension, and \( y \) is the location in the extra dimension.

Having setup the metric of the Randall-Sundrum model, the solution to the hierarchy problem is provided as follows. The model introduces the first brane called the “Planck” brane at \( y = 0 \), leaving the second brane called the “TeV” brane at \( y = r \). Because of the metric, mass scales are exponentially suppressed along the dimension \( y \), where masses \( (m) \) on the TeV scale brane are related to the mass \( (m_0) \) on the Planck scale brane: \( m = e^{-kr\pi} m_0 \). This way the weak scale is exponentially suppressed along the extra dimension, while the gravity scale is mostly independent of it. Since the values of all the bare parameters are determined by the Planck scale, an exponential hierarchy between the gravity scale and the weak scale can be naturally generated.

The Rundall-Sundrum model leads to a KK excitation to gravitational mode which would be observed in high energy collisions as a TeV-scale spin-2 Graviton \( (G^*) \). The bulk RS model is used in the measurements presented in this thesis as a benchmark model of neutral, spin-2 resonances that can decay to dibosons. It is characterized by a dimensionless coupling constant \( k/\bar{M}_{Pl} \), where \( k \) is the curvature of the warped extra dimension, and \( \bar{M}_{Pl} = M_{Pl}/\sqrt{8\pi} \) is the reduced Planck mass. This coupling constant is taken to be unity here in the analysis.

Information on the cross sections, widths, and branching ratios for bulk RS gravi-
tons at $\sqrt{s} = 8$ TeV are provided in Table 2.6. The cross sections are calculated at leading order from CalcHEP [22], using CTEQ6L1 PDFs and with the renormalization and factorization scales set to the graviton mass. Despite the recent availability of next-to-leading order predictions for RS models [23], the analysis is performed without K factors applied.

2.4.3 Heavy Vector Triplet

The heavy vector triplet (HVT) model [24] introduced a simplified phenomenological Lagrangian for a heavy vector triplet. This thesis interprets the combined diboson results in a model-independent way through HVT model.

Only relevant couplings and mass parameters are kept in the HVT model since the resonance searches are typically sensitive only to those parameters that control the resonance mass and the interactions involved in its production, not to all the details of the underlying model. The experimental results are expressed by confidence level
Table 2.6: Cross section, width, and branching ratios for bulk RS Graviton production at $\sqrt{s} = 8$ TeV.

<table>
<thead>
<tr>
<th>Pole Mass [TeV]</th>
<th>$\sigma$ [fb]</th>
<th>Width [GeV]</th>
<th>$B(G^{*} \rightarrow t\bar{t})$</th>
<th>$B(G^{*} \rightarrow WW)$</th>
<th>$B(G^{*} \rightarrow ZZ)$</th>
<th>$B(G^{*} \rightarrow HH)$</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.30</td>
<td>38530.</td>
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<td>0.6292</td>
<td>0.3583</td>
<td>0.00900</td>
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</tr>
<tr>
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<td>0.2718</td>
<td>0.04988</td>
</tr>
<tr>
<td>0.45</td>
<td>5364.</td>
<td>14.72</td>
<td>0.3237</td>
<td>0.3985</td>
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</tr>
<tr>
<td>0.50</td>
<td>3106.</td>
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<td>0.4180</td>
<td>0.3356</td>
<td>0.1795</td>
<td>0.06426</td>
</tr>
<tr>
<td>0.55</td>
<td>1870.</td>
<td>22.25</td>
<td>0.4805</td>
<td>0.2940</td>
<td>0.1556</td>
<td>0.06741</td>
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<tr>
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<td>0.65</td>
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<td>0.5773</td>
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</tr>
<tr>
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<td>0.5947</td>
<td>0.2180</td>
<td>0.1126</td>
<td>0.07270</td>
</tr>
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<td>0.80</td>
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<td>41.06</td>
<td>0.6083</td>
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<td>0.85</td>
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</tr>
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<td>0.6280</td>
<td>0.1958</td>
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</tr>
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curves in the parameter space of the phenomenological Lagrangian, easily translated into any specific model where the phenomenological parameters can be computed explicitly. Because of this simplification, the model cannot predict non-resonant properties such as off-shell production rates accurately.

The dynamics of the new vector is described by a simple phenomenological Lagrangian

\[
\mathcal{L}_V = -\frac{1}{4} D_{[\mu}V_{\nu]}D^{[\mu}V^{\nu]} + \frac{m^2}{2} V^a V^a \\
+ ig_V c_H V^a \tau^a \epsilon^\mu \mathcal{D}^\mu H + \frac{g^2}{2} c_F j_F^a \\
+ \frac{g_V}{2} c_{VVV} \epsilon_{abc} V^a V^b D^{[\mu}V^{\nu]} c + g^2 c_{VVHH} V^a V^a H^\dagger H - \frac{g^2}{2} c_{VVW} \epsilon_{abc} W^{\mu \nu \lambda} V^b V^c.
\]

If we look at the above equation, the first line contains the \( V \) kinetic and mass term, plus trilinear and quadrilinear interactions with the vector bosons from the covariant derivatives

\[
D_{[\mu}V_{\nu]} = D_\mu V_\nu - D_\nu V_\mu, \quad D_\mu V^a = \partial_\mu V^a + g \epsilon^{abc} W^b V^c,
\]

where \( g \) denotes the \( SU(2)_L \) gauge coupling.

The second line contains direct interactions of \( V \) with the Higgs current

\[
i H^\dagger \tau^a \epsilon^\mu \mathcal{D}^\mu H = i H^\dagger \tau^a D^\mu H - i D^\mu H^\dagger \tau^a H,
\]

and with the SM left-handed fermionic currents

\[
J_F^\mu = \sum_f \bar{f}_L \gamma^\mu \tau^a f_L,
\]

where \( \tau^a = \sigma^a / 2 \).

The third line of Eq. (2.16) contains 3 new operators and free parameters, \( c_{VVV} \), \( c_{VVHH} \) and \( c_{VVW} \). To a first approximation the operators in the third line can be
disregarded and the phenomenology can be described entirely by the four parameters $c_H$, $c_l$, $c_q$ and $c_3$, and the mass term $m_V$.

The new heavy vector triplet field $V^{\pm,0}$ couples to fermions by a combination of parameters $(g^2/g_V)c_F$ and to the Higgs and SM bosons via $g_Vc_H$. Here $g$ is the gauge coupling constant, $g_V$ represents the strength of the new vector bosons while $c_F$ and $c_H$ are the coupling strength to fermions and higgs which are close to be unity for most models. Two benchmark models with different values of $g_V$ are explored. One of the models (model A) allows the branching fractions to gauge bosons and fermions to be comparable, as in some extensions of the SM gauge group such as the EGM discussed previously. The other model (model B) strongly suppress couplings to fermions, as in the composite Higgs models. Figure 2.8 shows the allowed and excluded regions in the $(c_H, c_F)$ plane at 2 TeV. The points corresponding to models A and B for the different values of $g_V$ are also shown. $g_V = 1, 3, 6$ corresponds to weak, intermediate and strong coupling respectively. In the plots, the yellow region represents the exclusion from $V \to \ell\nu$ search [25], while the dark and light blue ones show the limits from $WZ \to 3\ell\nu$ [26] and $WZ \to jj$ [27] with $W/Z$ tagged jets respectively. The solid black lines depict constrains from electroweak precision tests at 95% CL and the dashed black line twice this limit.

Information on the cross sections, widths, and branching fractions for HVT production at $\sqrt{s} = 8$ TeV are provided in Table 2.7. The A($g_V=1$) scenario is studied here as a benchmark. It serves as a good reference point since it is very close to EGM $W'$. More information on the branching ratios for explicit models can be found in Ref. [24].
Figure 2.8: Current experimental constrains in the \((c_H, c_F)\) plane for the benchmark points at 2 TeV [24].

Table 2.7: Cross section (leading order), width, and branching ratios for HVT A \((gv=1)\) production at \(\sqrt{s} = 8\) TeV.

<table>
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<tr>
<th>Pole Mass [TeV]</th>
<th>(\sigma) [pb]</th>
<th>Width [GeV]</th>
<th>(B(V^0 \rightarrow WW))</th>
<th>(B(V^\pm \rightarrow WZ))</th>
<th>(B(V^0 \rightarrow ll))</th>
<th>(B(V^\pm \rightarrow l\ nu))</th>
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3.1 The Large Hadron Collider

The Large Hadron Collider is built by the European Organization for Nuclear Research (CERN), located across the border of France and Switzerland, near Geneva. It lies in a tunnel of approximately 27 km in circumference which is buried underground as deep as 175 m. LHC is designed to collide two opposing beams of protons ($p$) or heavy ions ($A$) at the central of energy of 13 TeV as of 2015. The design luminosity is $10^{34}$ cm$^{-2}$s$^{-1}$, corresponding to a bunch collision rate of 40 MHz.

The LHC is the last ring of the complex chain of the CERN accelerators, as shown in Figure 3.1. The proton sources are obtained from stripping electrons from the hydrogen gas by an electric field. The proton beams are then injected to a succession of machines–LINAC2, Proton Synchrotron Booster (PSB), Proton Synchrotron (PS), and Super Proton Synchrotron (SPS), which boost their energy to higher and higher levels until they reach their maximum energy of 4 TeV (as of 2012). Magnetic dipoles are used to bend the beam round via the Lorentz force while quadrupoles focus them.
to increase the chance of collision.

The two beams of protons travel counter to each other in separate pipes at close to the speed of light inside the accelerator. They are arranged to collide at four points along the ring where detectors record the collision data for research purposes—ATLAS [28], CMS [29], LHCb [30] and Alice [31]. Among them, ATLAS and CMS serve as the general purpose experiments for a large variety of physics analyses. More details about the ATLAS detector are presented in the following section.

![CERN's Accelerator Complex](image)

**Figure 3.1:** The CERN accelerator comlex showing the injectors used to fill the LHC and also some experiments located on the LHC ring [32].
3.2 The ATLAS Experiment

ATLAS (A Toroidal LHC ApparatuS) [28] is one of the two general purpose detectors at the LHC (located at Point 1), probing $p$-$p$ and $A$-$A$ collisions. It is sensitive to many interesting physics topics, such as high precision SM measurements, searches for new physics phenomena, and studies of the properties of the Higgs boson. A brief overview of the coordinate systems, the magnetic systems, the subdetectors as well as the trigger systems are presented below. A figure of the ATLAS detector is given in Figure 3.2 which also shows its important sub-systems.

![Figure 3.2: Schematic diagram of the ATLAS detector and its major sub-systems [28].](image)

3.2.1 Coordinate Systems

The origin of the right-handed coordinate systems lies at the nominal interaction point, while the $z$-axis is aligned along the beam direction and the $x$-$y$ plane is transverse to it. The $y$-axis is pointing to the center of the LHC ring and the $x$-axis is defined to point upwards. The azimuthal angle $\phi$ is measured from the positive
$x$ axis around the beam line, and the polar angle $\theta$ is the angle measured from the positive $z$ axis. The radial distance is given by $R = \sqrt{x^2 + y^2}$.

The rapidity is a very useful angular coordinate, defined as $y = \frac{1}{2} \ln \left( \frac{E+p_z}{E-p_z} \right)$, which is an invariant under a lorentz boost along the beam axis. It reduces to pseudorapidity $\eta = -\ln \left( \tan \left( \frac{\theta}{2} \right) \right)$ for massless particles. The angular distance $\Delta R$ in $\eta - \phi$ space is measured as $\sqrt{\Delta \eta^2 + \Delta \phi^2}$. Quantities in the transverse direction, i.e. the $x - y$ plane are very useful and common. For instance, the transverse momentum $p_T$ is defined as the momentum of the object transverse to the beam direction. Conservation of momentum can be easily applied given that the initial total transverse momentum is zero.

### 3.2.2 Inner Detector

The inner detector (ID) provides high resolution track measurements for charged particles over a pseudorapidity range of $|\eta| < 2.5$. It is surrounded by a magnetic field of 2 T provided by a solenoid which bends the charged particles in the ID as they leave hits in the active material of the detector. It consists of three independent but complementary sub-detectors: the pixel detector, the semiconductor tracker (SCT) and the transition radiation tracker (TRT), layered as shown in Figure 3.3. Figure 3.4 introduces the ID in a more technical way, showing not only each sub-system but also the geometric space that it covers.

#### 3.2.2.1 The Pixel Detector

The pixel detector is the inner-most part of the ID. It consists of 80 million readout channels, one for each pixel, providing position measurements as close as 5 cm from the beamline center. The system has three barrels along the beam (covering the central area) and three disks on each side (end-cap), which leads to three pixel layers to be crossed by each track. For the pixel barrel, the position resolution is 10 $\mu$m in
the transverse plane and 115 $\mu$m for the $z$ direction. For the end-caps, the intrinsic resolution is 10 $\mu$m in the transverse plane and 115 $\mu$m in the radial direction. Each pixel measures approximately 50 $\mu$m $\times$ 400 $\mu$m in $R-\phi$ and $z$ direction respectively. The nearest layer to the beam pipe is located at around 50 mm and covers |$\eta$| < 2.5.

3.2.2.2 Semiconductor Tracker

The next layer of the ID is the Semiconductor Tracker (SCT) which surrounds the pixel detector. The SCT contains four double layers of silicon strip detectors in the barrel region, and nine in the end-cap regions. The resolution in the barrel is 17 $\mu$m in the transverse plane and 580 $\mu$m in the $z$ direction, while the resolution in the end-caps is 17 $\mu$m in the transverse plane and 580 $\mu$m in the radial direction. Unlike the pixel detector, the innermost barrel modules are available up to |$\eta$| < 1.5, beyond which the end-cap modules are used. In total, the SCT has 6 million readout channels of silicon strips.
3.2.2.3 Transition Radiation Tracker

The third and the outermost part of the ID is the transition radiation tracker (TRT) which sits outside of the SCT. It is composed of 4 mm-diameter straws filled with a mixture of Xe (70%), CO$_2$ (23%) and O$_2$ (3%) gas interleaved with transition radiation material. The position resolution that the TRT provides is significantly worse than the Pixel or SCT; the only available information is on the $R-\phi$ direction for each hit with an accuracy of $\sim 130$ µm. The barrel is only available at $|\eta| < 1.0$, and the endcap covers up to $|\eta| \sim 2.0$.

Charged particles pass through the TRT and ionize the gas in the straws, therefore producing ionized electrons. The outer wall of the straw, covered by a thin layer of metal, is held at a high negative voltage which forces the ionized electrons to drift to the central anode wire. The time at which the free electrons get to the wire can be
used to determine how close the charged particle passes to the center of the straw. In addition, TRT is highly important for identifying electrons through the detection of transition radiation. Charged particles emit transition radiation photons when passing through TRT, and the probability of emitting a transition radiation photon is proportional to the Lorentz boost factor $\gamma$. Electrons are the only particles that have a large enough $\gamma$, thus emitting more transition radiation photons which can be used to distinguish electrons from other charged hadrons.

![Cut-away view of ATLAS calorimeter system](image)

**Figure 3.5**: Cut-away view of ATLAS calorimeter system [28]

### 3.2.3 Calorimetry

The particles reach the calorimeters after they make it through the ID. The calorimeter is designed to measure the energies of charged and neutral particles and to stop them before they reach the muon spectrometer. Muons and neutrinos are the only particles that are not strongly affected when passing through the calorimeters. An inner layer of electromagnetic calorimeter and an outer layer of hadronic calorimeter
form the ATLAS calorimetry. An overview of the ATLAS calorimeter is shown in Figure 3.5.

The calorimeter consists of metal plates (absorbers) and sensing elements. Interactions in the absorbers transform the incident energy into a “shower” of particles that are detected by the sensing elements. The part closest to the beam-line is a liquid argon calorimeter with lead as absorber and liquid argon as sensing elements. It consists of three cryostats, one barrel and two end-caps. The barrel cryostat contains the electromagnetic barrel calorimeter (EMB), whereas the two end-cap cryostats each contain an electromagnetic end-cap calorimeter (EMEC), a hadronic end-cap calorimeter (HEC), located behind the EMEC, and a forward calorimeter (FCal) to cover the region closest to the beam. The outer part is the hadronic calorimeter, with the sampling medium consisting of scintillator tiles and the absorber medium being steel. The hadronic calorimeter is extended to larger pseudorapidities by the HEC, a copper/liquid-argon detector, and the FCal, a copper-tungsten/liquid-argon detector. The hadronic calorimeter thus reaches one of its main design goals with a full coverage over $|\eta| < 4.9$.

A unique geometric design is used for both the barrel and end-cap sub-detectors, by which the various regions and layers are aligned in an accordion way. This geometry provides a perfect acceptance over the full $\phi$-symmetry range without leaving any gap. The accordion structure in use in the electromagnetic barrel can be seen in Figure 3.6.

3.2.4 Muon Systems

The muon spectrometer forms the outer part of the ATLAS detector which surrounds the calorimeter. It is designed to detect charged particles exiting the barrel and end-cap calorimeters and to measure the trajectories of them which are curved by the toroid magnetic field in the pseudorapidity range of $|\eta| < 2.7$. Over the range $|\eta| <$


\[ \eta = 0 \]

**Figure 3.6:** EM barrel module showing the three radial layers and granularity in \( \phi \) and \( \eta \) [28].

1.4, the muon tracks are bent by the large barrel toroid; over the range \( 1.6 < |\eta| < 2.7 \), the magnetic bend is provided by the two small magnets that are inserted to the two ends of the barrel toroid. Over the gap region \( 1.4 < |\eta| < 1.6 \), the magnetic deflection is produced by a combination of the barrel toroid and the end-cap field. The layout of the muon system is shown in Figure 3.7.

### 3.2.5 Trigger Systems

The trigger systems are designed to identify and select some of the most interesting events among a huge amount of input. Three levels of event selections–Level 1 (L1),
Figure 3.7: Cut-away view of ATLAS muon system [28]

Level 2 (L2), and Event Filter (EF) sub-systems apply stricter criteria to reduce the trigger rate sequentially. L2 and EF are known collectively as the High Level Trigger (HLT). L1 looks for events with signatures of high-$p_T$ electrons, photons, jets, $\tau$-leptons decaying to hadrons and large missing transverse energy. At L1, limited detector information such as coarse-grained calorimeter and muon information are combined to reduce the event rate to 75 kHz. Decisions are made within 2.5 $\mu$s and regions of Interest (RoI) are defined within the detector flagged by L1. L2 uses the full detector data within the RoIs at full granularity and precision and the rate can be decreased to roughly 3 kHz before the data has to be transferred from the full detector readout. The EF uses the offline analysis procedure and further reduces the event rate to about 200 Hz. With the help of the inner detector reconstruction information, the HLT algorithm provides the best particle identification.
3.3 Current Status for the ATLAS Detector and LHC

At the time of writing this thesis, the LHC Run 2 is ongoing at a center of mass energy of 13 TeV with proton bunch spacing of 25 ns. The estimation is that 100 fb$^{-1}$ of data will be collected by the time of the next shutdown in 2018. For LHC Run 2, the insertable B-layer (IBL) [33] is added as the fourth which is the innermost layer of the pixel detector. This has been one of the major upgrade activities performed during the LHC long shutdown 1. The IBL is designed to improve the precision of vertexing and $b$ tagging performance, therefore resulting in higher sensitivities in those analysis involving $b$ jets.
Monte Carlo Simulation of Proton Collisions

Simulated physics events using Monte Carlo (MC) methods provide a useful way to make predictions based on theoretical models and compare to data.

4.1 Monte Carlo Event Generation

It is extremely challenging to understand High Energy Physics (HEP) experimental data, for example at the LHC, to compare them with theories of particle physics. Data is an inclusive collection of all possible processes in the SM and beyond. It is not easy to compute the matrix elements beyond certain orders of perturbation theory, while QCD processes intrinsically include the non-perturbative feature of confinement which is not solved in the current state-of-the-art quantum field theory. There is also the problem of dealing with many divergences and/or near-divergences of the matrix elements, which cannot be easily solved in the theory.

It is, therefore, extremely important to make use of the technique developed in the past couple of decades, MC event generation, to help obtain reasonable theoretical predictions of experimental observables. Most MC tools rely on factorization, which separates the processes by momentum scale. At high scales, when high energy
exchange is involved, the energetic incoming partons produce a number of energetic outgoing partons, which is later referred as the “hard process”. It is the process where the matrix elements are perturbatively computable. At low scales, MC techniques describe soft processes: initial and final-state radiation, the parton shower, and the underlying event. Putting things together, the MC event generators are capable of simulating separately different physics processes of interest based on the different “hard processes”, on top of which “soft processes” are added.

From the experimental point of view, what a detector sees can only be hadrons rather than partons. Because of its non-perturbative nature, the hadronization process can only be modeled with tunable parameters to describe the hadron-level properties.

As part of modern high energy physics experiments, MC event generation performs an essential role in various aspects, such as the simulation of background processes, estimation of signal event rates and topologies, study of detector imperfections and so on. A general purpose generator is designed to mimic the whole process of producing and measuring actual collisions. The role of the MC generators is a virtual copy of a real machine, as presented in Figure 4.1.

The section is organized as follows. It starts with the introduction of how MC handles the “hard processes” in section 4.1.1. In section 4.1.2, it describes the treatment of the “parton shower” and “multiple parton interactions”. It is followed by the introduction to “hadronization” in section 4.1.3. The description of the “underlying event” is provided in section 4.1.4.

4.1.1 Hard Processes

One obvious appreciable feature of MC event generation is that it can simulate a typical process of interest, which can be quite rare in proton-proton collisions. Most the processes of interest in LHC have large momentum transfers, in which the
Figure 4.1: A work flow comparing real proton–proton collisions from LHC and virtual event simulation with MC generators.

involved partons are asymptotically free according to QCD and can be described by perturbation theory. They are therefore computable with Feynman diagrams and are referred as a “hard process” due to the large momentum transfer.

As essentially used in the MC event generation, the cross section of a $2 \to n$ process in proton–proton collisions can be formulated as in Equation 4.1.

$$
\sigma = \sum_{1,2} \int dx_1 dx_2 \int f_1^p(x_1, \mu_F) f_2^p(x_2, \mu_F) d\hat{\sigma}_{2\to n}(\mu_F, \mu_R),
$$

(4.1)

and,

$$
d\hat{\sigma}_{2\to n}(\mu_F, \mu_R) = d\Phi_n \times \frac{1}{2s} |M_{2\to n}|^2(\Phi_n; \mu_F, \mu_R)
$$

(4.2)

where,

- $f_i^p(x_i, \mu_F)$ ($i = 1, 2$): PDFs for parton $i$, which depend on the momentum...
fraction of the parton to the parent proton, $x$, and factorization scale, $\mu_F$.

- $\hat{\sigma}_{2\to n}(\mu_F, \mu_R)$: parton level cross section with 2 partons in the initial state from the two colliding protons and $n$ partons in the final state. $\mu_R$ represents the renormalization scale.

- $\mathcal{M}_{2\to n}$: the tree level matrix element, which sums over relevant Feynman diagrams. The initial state spin and color degrees of freedom are averaged over, while the final state ones are integrated.

- $d\Phi_n$: differential phase space over $n$ final state partons. There are cuts involved to take into account the acceptance of the detectors, and to cut off the region where an emitted parton is becoming soft or collinear.

The cross section in Equation 4.1 of the hard process is valid only when the choices of factorization and renormalization scales and PDFs are fixed. There is no existing “correct” $\mu_F, \mu_R$ from first principles. However, our knowledge of the theory of QCD provides reasonable ranges. For example, in the process $2 \to 2$ or $2 \to 1$ of producing an s-channel resonance of mass $M$, the scales are often chosen as $\mu_F = \mu_R = M^2$. Regarding the choice of PDFs, each MC generator used in ATLAS uses its own default PDF, while at the same time provides access to other widely used PDF sets via the LHAPDF interface [34].

4.1.2 Parton Shower

Once the hard process matrix elements are computed to some fixed order of accuracy, the effects of the higher order terms are simulated with parton shower process. It includes a momentum transfer scale down to as low as 1 GeV, the level where confinement of partons into hadrons takes place. Similar to an electron emitting photons when moving in a field, quarks carrying colors can emit gluons. Unlike elec-
tron bremsstrahlung in QED, gluons also carry color charge so that they can emit gluons as well. This makes the parton shower even more complicated.

A most straightforward case to consider is that one parton splits into two partons, i.e. $q \rightarrow g + q$. The relationship of the differential cross sections before ($d\sigma_n$) and after ($d\sigma_{n+1}$) the splitting is characterized in Equation 4.3 [35]:

$$d\sigma_{n+1} = d\sigma_n \frac{\alpha_s}{2\pi} \frac{d\theta^2}{\theta^2}dzd\phi P_{ji}(z, \phi),$$  

(4.3)

where, the splitting process is $i \rightarrow j + k$. $\alpha_s$ is the strong interaction coupling constant. $\theta$ is the opening angle, while $\phi$ is the azimuthal angle of the splitting, from direction of parton $j$ to the direction of parton $i$, and $z$ is the energy fraction of parton $j$ to parton $i$. $P_{ji}$ is the $i \rightarrow j + k$ splitting probability function, which describes the distribution of the fraction $z$ of the energy of $i$ carried by $j$.

A parton shower is developed based on the sequential application of Equation 4.3 for each colored parton from the hard process. The evolution variable is taken as the virtual mass squared (virtuality, $q^2$) of the parton in the shower, based on which the parton shower is ordered. It starts from the $q^2$ close to, but lower than, the momentum transfer scale in the hard process, $q^2 < Q^2$. It falls when the shower process goes on because of the energy loss in the splitting. The shower process ends when $q^2$ falls to the hadronization scale around 1 GeV.

The application of Equation 4.3 only takes into account the collinear real parton emissions. The virtual effects, quantum loops, are included in the probability of not emitting partons during the evolution process i.e. from scale $q_1^2$ to $q_2^2$, characterized by the Sudakov form factor [35].

Besides the parton shower from final-state hard process partons as described above, a slightly different approach with same principle of parton emission is applied for parton radiation from initial-state hard incoming partons.
4.1.3 Hadronization

The partonic final state presented in the previous sections are still in the theory part of the Monte Carlo event generation. In order to compare with real proton–proton collisions in an experiment, one has to turn the partonic final state into a real hadronic final state which has been “seen” by the detectors. The term “hadronization” stands for the model used in the Monte Carlo event generator.

There are two main hadronization models in use, the string model and the cluster model [36]. The former transforms the final-state partons directly into hadrons, while the latter makes uses of intermediate cluster objects with a mass scale of a few GeV.

The Lund string model [37] is the most well-known string model, which treats all but the highest-energy gluons as field lines. They are attracted to each other due to the gluon self-interaction, forming a narrow string of strong color field. However, in the string model, it is assumed that the fragmentation of a single string is in isolation, which may not be the case in proton-proton collisions. In this case, multiple parton interactions would lead to strings overlapping in space and time, especially at high energies.

On the other hand, the cluster model is based on the so called pre-confinement property of parton showers. In the so called Field-Fox-Wolfram model [38], the non-perturbative splitting of a gluon into a quark-anti-quark pair is enforced at the end of the parton shower evolution at a scale near a cutoff of $Q_0 \sim 1\text{GeV}$. These quark-anti-quark pairs are clustered into corresponding color pairs. Once the clusters are formed, the ones with mass $\lesssim 3\text{ GeV}$ can be decayed into hadrons using a simple isotropic two-body phase space model. The heavier mass cluster need to first go through the non-perturbative split into lower mass clusters before the cluster-to-hadron transition is performed. The process is repeated until no clusters are left untransformed.

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The cluster hadronization model is widely used in MC generators, i.e., HERWIG++ [39] and SHERPA [40].

4.1.4 Underlying Event

An “underlying event” (UE) includes the extra activities not identified with the given hard process in an event. Such activities include multiple parton interactions (MPI), QCD color connections between partons and beam remnants, as well as initial state radiation (ISR) and final state radiation (FSR).

The low momentum exchange partonic QCD process in the same collision event is the dominant contribution to UE due to its large production cross section. On the other hand, the soft interactions are not valid in perturbative theory, which makes it quite important to model them in MC. The related parameters are not known a priori, so they have to be tuned to match collision data.

In hadron colliders like LHC, the cross section of the partonic two-to-two scattering (e.g. $gg \rightarrow gg$) is huge, larger than the total proton–proton collision cross section. It means that on average more than one partonic collision happens in one proton–proton collision. The ratio of the two cross sections, $\sigma_{\text{partonic}} / \sigma_{\text{hadronic}}$, is interpreted as the mean number of partonic collisions in the event. The tuning of this variable for LHC collisions is discussed in Section 4.4.

4.2 Detector Simulation

GEANT4 (GEometry ANd Tracking) [41] [42] is a generic toolkit to simulate particles traversing the ATLAS detector. The configuration of the detector, including active detector elements, electronics, misalignment and known hardware issues, are specified in databases and libraries. This complete description of the detector is used to mimic the conditions of true data for comparison.

As a particle traverses the detector, energy deposit information is stored in a hit
file, which contains all input necessary for the next stage—digitization step as well as original information about the particles from event generation. The four-vectors of the particles are also recorded and set aside for later use.

There are also certain fast simulation processes available for quickly generating large MC samples, which do not rely on full Geant4 simulation.

4.3 Digitization

The energy deposits recorded in the “hits” need to be translated into the output actually produced by the ATLAS detector. The digitization step converts the simulated energy deposits inside sensitive detector material into the detector response, making them ready to be used for reconstruction. The digitization of Geant4 hits in ATLAS is performed subdetector-by-subdetector. Unlike the previous steps in the simulation chain, this is a very detector-specific task. Each subdetector has its own software to do the digitization, overlaying those effects not due to the primary collisions (pileup, beam halo, beam gas etc.). The output format is RAW Data Objects (RDOs). In addition to RDOs, Simulation Data Objects (SDOs) are created to save some simulation information which may be useful.

4.4 ATLAS Run 1 Pythia8 Tunes

4.4.1 Introduction to Tuning

As described before, experimental measurements rely on MC generators. Many models, such as hadronization, the parton shower, and multiple interaction are used to produce simulated events as similar as possible to observed data. Those models have \textit{a priori} unknown free parameters that need to be optimized to give an accurate description of the observables. However, most of the time they are correlated in a high dimensional parameter space. The process of the optimization of the parameters is known as tuning, and the parameter sets that are tuned are called MC generator
ATLAS used early LHC Run 1 data to tune the MPI parameters in Pythia 8 [43], generating the extensively-used AU2 tune for simulation of UE in hard interaction events, and the A2 tune for minimum bias (MB)/pileup event simulation [44].

My contribution to ATLAS generator tuning involved tuning the parton shower and multiparton interactions parameters of the MC generator Pythia 8 in a single step for a variety of observables based on ATLAS Run 1 data.

This study used two analysis toolkits: the Rivet generator-diagnostic validation system [45] was used to produce distributions of observables from a MC generator. The Professor tuning system [46] parameterized the results of MC tunes, and finding an optimal set of parameters given the data.

4.4.2 Tuning Setup

The tuned parameters are listed in Table 4.1 with a very brief definition, along with their sampling range used for the current tune. For example, if the final state of the hard process (not counting subsequent resonance decays) contains at least one quark, gluon or photon, then \( p_T^{\text{max}} \), which is the maximum \( p_T \) to be allowed in the shower evolution, is chosen to be the factorization scale for internal processes referred as \( p_T^{\text{factorization}} \). The parameter \( p_T^{\text{maxFudge}} \) introduces a multiplicative factor \( f \) such that \( p_T^{\text{max}} = f \cdot p_T^{\text{factorization}} \), and this parameter only applies to the hardest interaction in an event. \( f \) is often set to 1, but variations around this default can be useful as tests. More explanation of the meanings of the parameters can be found in [47].

The observables used for the tuning are taken from a series of analyses which are listed in Table 4.2. The measured and predicted observables are compared in bins, and the predicted contents of each bin are parametrized by a multi-dimensional 3rd order polynomial in the parameter hyperspace, based on 500 sampled tune points. Then they are used to calculate a \( \chi^2 \) with respect to the reference data, which
allows us to numerically determine the best tune point in the parameter space which corresponds to the point. Weight factors are added in the $\chi^2$ to place emphasis on certain observables.

### 4.4.3 **POWHEG+PYTHIA8 Shower Tuning Submission Framework**

As there is the need for streamlined infrastructure to submit and validate Monte Carlo production, a shower tuning submission framework has been developed. I created a framework based on the standalone package Sacrifice which steers PYTHIA8 without Athena. One of the features of the framework is that it provides interfaces to different hard-scatter generator output formats (LHEF\textsuperscript{1}[48] and HepMC\textsuperscript{2}[49]) as well as different PDFs (LHAPDF\textsuperscript{3}[34]). It also reads existing LHEF files for analysis needed, then submits jobs to the ATLAS computing grid, which is a collection of computer resources from multiple locations. The framework makes it easy to store and retrieve data from distributed storage systems and to monitor all of these operations, and is widely used across the MC group of ATLAS collaboration.

---

\textsuperscript{1} the Les Houches Event File

\textsuperscript{2} a C++ event record for MC generators

\textsuperscript{3} the Les Houches Accord PDFs
Table 4.2: Summary of tuning analysis

<table>
<thead>
<tr>
<th>Tuning analysis number</th>
<th>Note</th>
</tr>
</thead>
<tbody>
<tr>
<td>ATLAS-2011-S8924791</td>
<td>Jet shapes at 7 TeV in ATLAS</td>
</tr>
<tr>
<td>ATLAS-2011-S8971293</td>
<td>Dijet azimuthal decorrelations</td>
</tr>
<tr>
<td>ATLAS-2011-I919017</td>
<td>Measurement of ATLAS track jet properties at 7 TeV</td>
</tr>
<tr>
<td>ATLAS-2011-S9126244</td>
<td>Measurement of dijet production with a veto on additional central jet activity</td>
</tr>
<tr>
<td>ATLAS-2012-I1119557</td>
<td>More Jet shapes</td>
</tr>
<tr>
<td>ATLAS-2012-I1094564</td>
<td>Jet mass and substructure of inclusive jets</td>
</tr>
<tr>
<td>ATLAS-2012-I1094568</td>
<td>ttbar jet veto</td>
</tr>
<tr>
<td>ATLAS-2011-S9131140</td>
<td>Measurement of the Z $p_T$ with electrons and muons</td>
</tr>
<tr>
<td>ATLAS-2012-I1084540</td>
<td>Rapidity gap cross sections measured with the ATLAS detector in pp collisions at 7 TeV</td>
</tr>
<tr>
<td>ATLAS-2012-I1183818</td>
<td>Forward ET flow</td>
</tr>
<tr>
<td>ATLAS-2011-S8994773</td>
<td>Calo-based underlying event at 900 GeV and 7 TeV in ATLAS</td>
</tr>
<tr>
<td>ATLAS-2010-S8918562</td>
<td>Track-based minimum bias at 900 GeV and 7 TeV in ATLAS</td>
</tr>
<tr>
<td>ATLAS-2010-S8894728</td>
<td>Track-based underlying event at 900 GeV and 7 TeV in ATLAS</td>
</tr>
<tr>
<td>MC-XS</td>
<td>MC analysis for process total cross section</td>
</tr>
</tbody>
</table>

4.4.4 Results

Table 4.3 summarizes four sets of tuning parameters and their comparisons. The default value of the tune is set as a reference tune; tune1 is obtained by dynamically varying the tuning parameters to get the best performance; tune2 is obtained by setting the parameter ISR $\alpha_s$ fixed while floating the others; tune3 is obtained by setting both the parameter ISR $\alpha_s$ and the parameter FSR $\alpha_s$ fixed while floating the others. Default, tune1, tune2 and tune3 correspond to the magenta, green, blue and red lines separately in Figure 4.2 to Figure 4.5. Definitions of the variables in the plots can be roughly explained as follows:

- $\Psi(r), \rho(r)$ [50]: The integrated jet shape $\Psi(r)$ is defined as the average fraction
of the jet $p_T$ that lies inside a cone of radius $r \ (r = \sqrt{\Delta y^2 + \Delta \phi^2})$ concentric with the jet cone. The differential jet shape $\rho(r)$ is defined as the average fraction of the jet $p_T$ that lies inside an annulus of inner radius $r - \Delta r/2$ and outer radius $r + \Delta r/2$ around the jet axis.

- Gap fraction, $Q_0$ [51]: Gap fraction is the fraction of events that do not have an additional jet with a transverse momentum greater than a given veto scale, $Q_0$, in the rapidity interval bounded by the dijet system.

- $\tau_{21}, \tau_{32}$ [52][53]: The N-subjettiness variable $\tau_N$ effectively "counts" the number of subjets in a given jet. $\tau_2/\tau_1$ ($\tau_{21}$) is found to be an effective discriminating variable to identify two-prong objects, such as boosted $W$, $Z$, and Higgs bosons and $\tau_3/\tau_2$ ($\tau_{32}$) is effective for three-prong objects like boosted top quarks.

- $\sqrt{d_{12}}, \sqrt{d_{23}}$ [54]: The $k_T$ splitting variable is defined as $\sqrt{d_{ij}} = \min(p_{T,i}, p_{T,j}) \times \delta R_{i,j}$, where $i$ and $j$ represent the two proto-jets combined at the final ($\sqrt{d_{12}}$) or second-to-last ($\sqrt{d_{23}}$) step of the $k_T$ algorithm.

For observables that measure the fraction of events passing a veto requirement in events with a dijet topology, a significant improvement can be seen in the new tunes as opposed to the reference tune. The description of the jet substructure observables is affected by FSR, and is almost all within the uncertainty bands. Slight improvement can be seen for $Z p_T$ and jet shapes. The application of these tunes for ATLAS sample production is intended to be mainly used for situations where next-to-leading order Monte Carlo is unavailable.
Figure 4.2: Results of Jet Shape analysis comparing different tunes. The yellow shaded area represents the data uncertainty.

Figure 4.3: Results of $Z p_T$ measured using electron and muon channels comparing different tunes. The yellow shaded area represents the data uncertainty.
Figure 4.4: Results from dijet events measuring the fraction of events passing a veto requirement comparing different tunes. The yellow shaded area represents the data uncertainty.

Table 4.3: Summary of tuning parameters’ comparisons. Tune1 represents reference tune; tune2 is fixing FSR $\alpha_S$; tune3 is fixing both ISR $\alpha_S$ and FSR $\alpha_S$.

<table>
<thead>
<tr>
<th>Tuning parameters</th>
<th>Default value</th>
<th>tune1</th>
<th>tune2</th>
<th>tune3</th>
</tr>
</thead>
<tbody>
<tr>
<td>SpaceShower:alphaSvalue</td>
<td>0.137</td>
<td>0.1258525</td>
<td>0.1399655</td>
<td>0.125</td>
</tr>
<tr>
<td>TimeShower:alphaSvalue</td>
<td>0.1383</td>
<td>0.1505813</td>
<td>0.125</td>
<td>0.125</td>
</tr>
<tr>
<td>BeamRemnants:primordialKThard</td>
<td>2.0</td>
<td>1.631222</td>
<td>0.162261</td>
<td>1.841699</td>
</tr>
<tr>
<td>SpaceShower:pTmaxFudge</td>
<td>1.0</td>
<td>0.8936378</td>
<td>0.8287265</td>
<td>0.8473</td>
</tr>
<tr>
<td>SpaceShower:pT0Ref</td>
<td>2.0</td>
<td>1.609147</td>
<td>0.1853912</td>
<td>1.78297</td>
</tr>
</tbody>
</table>
Figure 4.5: Results from jet substructure analysis comparing different tunes. The yellow shaded area stands for data uncertainty.
This chapter summarizes the particle identification and the object reconstruction in the ATLAS detector. Figure 5.1 shows how the particles are detected by the ATLAS detector.

5.1 Inner Detector Tracks

Inner detector tracks are very important for reconstructing electrons and muons, as well as determining the primary vertices and which particles originate from them. Tracks are reconstructed using two algorithms: the inside-out sequence and the outside-in sequence [56].

The inside-out reconstruction sequence starts from a collection of hits in the pixel detector and the SCT. Track candidates are formed and seeded from these hits and are propagated. Additional hits in the detector elements are added along the track road through accepting and rejecting, with the track information updated as each hit is added. Ambiguities between seeds must be resolved in the meanwhile. All tracks are ranked by a scoring system which favors long tracks over short segments and takes into account the fit quality for each track. Hits associated with more than one
track are assigned to whichever track that has a higher score, and the score of that track is updated instantaneously. Finally, after all hits are assigned to some track, those tracks with scores below a threshold are discarded. The tracks are extended to TRT and matched with TRT hits to form the complete inner tracks.

The outside-in reconstruction sequence is useful when tracks cannot be seeded by the inside-out sequence. This can occur in several cases. For example, a hadron decay that happens outside of the pixel detector is unlikely to be seeded; electrons coming from the photon conversions do not leave hits in the silicon detector and are unlikely to be seeded either. In this case, stand-alone TRT tracks are formed and worked backwards into the silicon detector. Due to the lack of resolution in the straw direction, this sequence works with low efficiency and serves mainly to correct for
the intrinsic sources of inefficiency of the inside-out sequence.

Primary vertices are found by using an iterative vertex finding algorithm [57]. To define a vertex, a minimum of two tracks are required. The vertex with highest scalar $p_T$ sum of tracks is called the hard scatter vertex, and any physics objects in the analyses are usually required to originate from it.

5.2 Calorimeter Clusters

Calorimeter clusters are constructed from energy deposits in cells in the electromagnetic and hadronic calorimeters. A topological clustering algorithm [58] is used to group neighboring cells together into clusters, which are mainly used for jet reconstruction in the following analyses. The algorithm proceeds in two steps: cluster formation and cluster splitter. First, a list of seeds is formed in descending order with signal to noise ratio $S/N$ above some threshold $t_{seed}$. Then any neighboring cell to the seed with a $S/N$ above some threshold $t_{neighbor}$ which is not also a seed is added to form a proto-cluster. If a neighboring cell is close to more than one proto-cluster, the proto-clusters are merged. The procedure is repeated until all the neighboring cells are added to some proto-cluster. A cut $E_T > E^{threshold}_T$ is applied to the list to form the final proto-cluster and the cluster formation step is complete. The cluster splitter step repeats the same procedure, but the seed cells are defined as local maxima with $E > 500$ GeV and only those cells in the proto-clusters from the first round are involved. The proto-clusters formed in this step are the final topoclusters. Different threshold values are set for the EM calorimeter and the hadronic calorimeter. For the 8 TeV collider run, in the EM calorimeter, $t_{seed} = 6$, $t_{neighbor} = 3$, and $E^{threshold}_T = 5$ GeV; in the hadronic calorimeter, $t_{seed} = 4$, $t_{neighbor} = 2$, and $E^{threshold}_T = 0$. 

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5.3 Electrons

An electron is defined as one or more tracks in the ID matched with an energy cluster in the electromagnetic calorimeters in the precision measurement region of the detector, $|\eta| < 2.5$. The standard electron reconstruction starts with seed clusters from the sliding-window algorithm [58]. Tracks are matched to seed clusters by extrapolating from the last measurement point to the center of the EM calorimeter. A track is considered to be a match if it is within a distance of $\Delta \eta < 0.05$ and $\Delta \phi < 0.1$. After the matching procedure, the cluster energy is re-computed using the process discussed in Section 5.3. The final cluster energy is assigned as the electron four-vector energy, while the $\eta$ and $\phi$ directions of the electron are decided based on the matched track. Other algorithms for reconstructing low $p_T$ electrons and forward electrons (with $|\eta| > 2.5$) are not used in the analyses in this thesis and will not be discussed here.

The identification for standard electrons is based on the variables of tracks and clusters. The analyses discussed in this thesis use a cut-based electron identification selection. The $\text{loose}^{++}$, $\text{medium}^{++}$, and $\text{tight}^{++}$ working points are chosen to reject fake electrons from photon conversions, Dalitz decays, and light jets. In order to reject the electrons from photon conversion, at least one hit in the innermost pixel detector layer (the b-layer) is required. Further requirement on the ratio of the cluster energy and the track momentum is also applied. Additional cuts, for example isolation cuts, are often imposed on electrons after identification requirements to remove electrons with significant nearby hadronic activity. The electron track isolation discriminant is computed as the summed scalar $p_T$ of additional tracks inside a cone of radius $\Delta R$ around the electron. The tracks considered must originate from the same vertex associated with the electron. The electron calorimeter isolation discriminant is defined as the scalar sum of the transverse energy $E_T$ of topological
clusters within a cone of radius $\Delta R$ around the electron cluster and is corrected for any contribution from the electron energy and pileup.

5.4 Muons

Muons are reconstructed using information from the muon spectrometer (MS) which can be matched with ID tracks. There are a variety of strategies to reconstruct and identify muons in ATLAS: stand-alone, combined, and tagged muons.

Stand-alone muons are formed from tracks in the muon spectrometer, extrapolated back to the interaction point. Combined muons are seeded by stand-alone muon tracks and are combined with inner detector tracks. Segment-tagged muons are formed by extrapolating inner detector tracks through the first station of the muon system and searching for nearby track segments.

For use in analyses, these three algorithms can be combined to increase the muon reconstruction efficiency. In the combination process, if a stand-alone muon is combined with more than one ID track, a best match inner detector track is selected based on $\eta - \phi$ matching. In addition, an overlap removal between the algorithms is applied. For instance, a stand-alone muon will be removed if it uses the same best match ID track as a combined muon. The muon track and calorimeter isolation discriminants are the same as the ones used for electrons, except for the isolation cone radius can be slightly different. More details can be found in the ATLAS muon reconstruction conference note [59].

5.5 Jets

Jets are collimated streams of hadrons and are essentially the key observables for the following analyses. Individual jets are proxies for quark and gluons, and can be defined by particle clustering algorithms as introduced in Section 2.3.
5.5.1 Calorimeter Jets

The calorimeter system, made up of the EM calorimeter and the hadronic calorimeters are used as the principal detectors for jet reconstruction. In principle, the pseudorapidity coverage of the detectors is near hermetic, $|\eta| < 4.9$. Several jet algorithms in ATLAS are provided to meet different purposes of physics analysis as described in Chapter 2.

The jet energy is corrected for non-compensation and inhomogeneities of the calorimeter with the help of the calibration schemes. The energy resolution for the central barrel is $\sigma(E)/E = 10\% + 0.7\%$ which impacts the width of resonances reconstructed with jets, and the granularity is around $\Delta \eta \times \Delta \phi = 0.025 \times 0.025$ in the 2nd layer of the EM calorimeter. This fine granularity determines the smallest scale at which substructure can be measured in the calorimeter.

5.5.2 Track Jets

Track jets are constructed with the same clustering algorithms as standard jets, using the ID tracks as the inputs. Besides the calorimeter jets, jet finding from the ID tracks can be useful since it can improve the quality of the jet parameters. ID tracks could also identify the reconstructed vertex associated with the jet. The double ratio of the track jet and calorimeter jet transverse momenta in data and MC is used for the validation of the jet energy scale uncertainty.

5.5.3 Ghost Association of Tracks to Jets

One of the important ways to avoid the effect of pileup (additional $p$-$p$ interactions produced within the same event) is to use objects that are inherently immune to pileup, such as tracks. The very good $z$ resolution of tracks allows efficient matching of tracks to vertices, and rejection of pileup tracks. Then jets can be classified according to the fraction of non-pileup tracks they contain. As a result, we can
take advantage of the tracking information when constructing the jets by associating tracks to jet. Traditionally, tracks were associated to jets by geometrical matching ($\Delta R$ cone), however jets are not always circular.

Ghost association [60] is a brilliant method of associating tracks to jets by involving them in the clustering. Representatives of the tracks with infinitesimally small energy and zero mass, called ghosts, are added to the list of constituents for the clustering of forming jets, assuming the $\eta$ and $\phi$ coordinates remain the same. Then the jet clustering procedure is repeated, which should be able to produce the same set of jets as before\(^1\). That is because any valid jet clustering algorithm would not be affected by adding a tiny object. By this method, each track is associated with a unique jet.

Because a jet’s ghost-associated tracks from the primary vertex is insensitive to pileup, this quantity is used as a tagging variable in the the following analyses.

5.5.4 Jet Calibration

The calorimeter cells measure energy at the electromagnetic scale, giving the correct energy for an electromagnetic shower deposit. However, hadrons have a lower response which is not accounted for, as well as some energy escaping due to dead material in the detector, or not being detected above the level of noise. As a result the jets are calibrated to provide a more accurate measurement of the energy. Topoclusters, introduced in Section 5.3, are calibrated using the local cluster weighting (LCW) calibration scheme [61] before building the jets. Once all clusters have been calibrated, they are grouped into jets using the jet algorithms discussed in Chapter 2. A further calibration—the jet energy scale (JES) calibration, derived from the response of jets built from final state truth electrons, photons, and hadrons, is ap-

\(^1\) Depending on the choice of $\beta$ in the formulae 2.13 in Chapter 2, the jets can be affected by adding a large number of ghosts at once.
plied to jets. For the analyses discussed in later chapters, a jet calibration scheme is derived based on QCD MC jets, resulting in LCW+JES jets. One of the aims of this calibration is to correct for the non-uniformity of the jet energy and mass response as a function of the jet direction in the detector.

The jet response is defined as

\[ R_{\text{jet}}^E(\eta_{\text{det}}) = \frac{E_{\text{jet}}^{\text{reco}}}{E_{\text{jet}}^{\text{truth}}}, \quad R_{\text{jet}}^m(\eta_{\text{det}}) = \frac{m_{\text{jet}}^{\text{reco}}}{m_{\text{jet}}^{\text{truth}}}. \]

The LCW calibrated jet is corrected as follows:

\[ E_{\text{jet}}^{\text{LCW+JES}} = \frac{E_{\text{jet}}^{\text{LCW}}}{\mathcal{F}_{\text{calib}}(E_{\text{jet}}^{\text{LCW}})}, \quad m_{\text{jet}}^{\text{LCW+JES}} = \frac{m_{\text{jet}}^{\text{LCW}}}{\mathcal{F}_{\text{calib}}(E_{\text{jet}}^{\text{LCW}} \ast \text{JES})}, \]

where \( E_{\text{jet}}^{\text{LCW}} \) and \( m_{\text{jet}}^{\text{LCW}} \) are the energy and mass of the jet at the LCW scale, respectively. \( \mathcal{F}_{\text{calib}}(E_{\text{jet}}^{\text{LCW}}) \) is the calibration function, which is derived in different \( \eta_{\text{det}} \) and \( E_{\text{truth}} \) bins. For each \( \eta_{\text{det}} \) and \( E_{\text{truth}} \) bin, the mean jet energy (mass) response \( R_{\text{jet}}^E(\eta_{\text{det}}) \) (\( R_{\text{jet}}^m(\eta_{\text{det}}) \)) is derived as the peak position of the Gaussian fit to the response distribution.

Figure 5.2 illustrates how calibration works. Before the energy correction, a clear shift of the energy response as well as a detector dependence is visible. After the jet energy correction, the energy response is flat and centered around its expected value.

5.5.5 Jet Quality Requirements

The jets not associated to in-time real energy deposits in the calorimeters are tagged as “bad” jets. Certain problems such as hardware errors, poor beam conditions and cosmic ray showers can cause a jet to be flagged as a “bad” jet. Several variables concerning the jet energy deposition fraction in different part of the detectors, jet timing and jet quality are used to identify such a “bad” jet. If a “bad” jet is found in the event, the entire event will be regarded.
Figure 5.2: Jet energy response before (top) and after (bottom) jet energy calibration as a function of the jet $\eta$ for different energies [62].
5.6 Missing Transverse Energy

The missing transverse energy ($E_T^{\text{miss}}$) is calculated to estimate the summed transverse momentum of all particles produced in the collisions that hardly interact with the detectors, such as neutrinos. It is defined as the negative vector sum of all particle momenta in the plane transverse to the beam, composed of two terms—the calorimeter and muon terms. The calorimeter term includes all energy deposits in the calorimeter, such as electrons, photons, jets, $\tau$ leptons. Each of these objects is calibrated to its final energy scale to improve the resolution of the $E_T^{\text{miss}}$. The muon term is calculated as the total transverse momentum of all identified muons.
A search is performed for diboson resonances (WW, WZ and ZZ) in fully hadronic final states. The massive resonances result in W or Z bosons, and their hadronic decay products are each reconstructed as a single jet (also known as VVJJ analysis, where V denotes the W or Z boson and J denotes the jet). The dominant background is dijet events from QCD processes, which provides a smoothly falling spectrum. Thus the goal is to look for resonance structures in the dijet invariant mass spectrum. To suppress the dijet background, boson tagging techniques and jet substructure variables are used to select jets. Limits on the production cross section times branching ratio of massive resonances are set in each diboson channel as a function of the resonance mass using the benchmark models. This chapter describes in detail the analysis techniques used to look for new resonances.
6.1 Data Samples

The data used in the analysis was recorded in 2012 at a center of mass energy of 8 TeV corresponding to a luminosity of 20.3 fb\(^{-1}\). The events are selected based on a single-jet trigger with a \(p_T\) threshold of 360 GeV. They are also required to pass some of the recommended data quality cuts. A first requirement of data quality is a selection called the good run list (GRL), which ensures that each sub-detector system has its performance inspected for each period and that the system is in good working order.

6.2 Monte Carlo Samples

The process \(W' \rightarrow WZ\) is used as a benchmark process for this analysis. The resonance decays to a \(W\) and a \(Z\) boson, which are then both forced to decay hadronically. \(W'\) MC samples are produced with PYTHIA8 between 1.3 and 3.0 TeV at 100 GeV intervals. These are based on the EGM model. An additional set of \(W'\) samples generated with PYTHIA8 for the hard scattering interaction and HERWIG++ for parton showering and hadronisation is used to help in the assessment of systematic uncertainties.

The processes \(G^* \rightarrow WW\) and \(G^* \rightarrow ZZ\) are also used as benchmark processes. The resonance decays to \(WW\) or \(ZZ\), which are then forced to decay hadronically. The coupling constant \(k/\bar{M}_{Pl} \sim 1\), as described in Section 2.4.2. The samples are produced with CalcHEP 3.4 [22] setting \(k/\bar{M}_{Pl} = 1\) between 1.3 and 3.0 TeV at 100 GeV intervals. Events are generated using CTEQ6L1 [63] PDFs, and use PYTHIA8 for the parton shower and hadronization.

Main backgrounds for the analysis are QCD multijet processes, \(V+\)jets and diboson events. The QCD background is by far the most dominant. Monte Carlo samples generated with PYTHIA8 are used to assess both the QCD and \(V+\)jets backgrounds,
the latter of which are found to be negligible. Diboson events are produced at the
generator level with Powheg, using Pythia for the soft parton shower.

6.3 Jet Grooming and Mass Drop Filtering

The jet grooming procedure aims to identify concentrations of energy within a jet and
remove the soft structure in a jet. The hope is that constituents from the underlying
event or pileup can be discarded by the grooming step.

Mass-drop filtering is a jet grooming technique which attempts to find the hard
substructure within a jet. It turns out to be very useful for analysis that involves
boson decay. In this thesis, a mass-drop filtering algorithm called BDRS-A which is a
modification of the BDRS mass-drop filtering algorithm [64] is used. BDRS-A uses a
geometric distance measure algorithm, such as the Cambridge-Aachen algorithm, to
reconstruct the original jet rather than an algorithm with a distance measure based
on $p_T$. Because this reconstruction procedure is only based on angular distances,
large splittings in structure within a jet can be identified while the actual angular
scale (which varies with the boost of a decaying boson) remains flexible. The idea
of the BDRS-A algorithm is to determine whether a gluon is emitted or not in the
large-R jet, rejecting these massive QCD jets by measuring the energy fraction of
a splitting. Given a large-$R$ jet, we apply the following iterative decomposition
procedure—the mass-drop technique: First, split the large-$R$ jet $j$ into two subjets
by undoing the last step of the C/A algorithm and labeling the one with the greater
mass as $j_1$ and the other as $j_2$. If there exists a mass difference (parameterized by
a mass ratio parameter) between the original jet $j$ and $j_1$ and the splitting in $p_T$ is
rather symmetric, then exit the loop and let $j$ to be the final jet. Otherwise, discard
$j_2$ and redefine $j$ to be $j_1$. This way the subjet with lower mass is removed at each
step, and the jet with higher mass is split. This split-and-discard procedure keeps
going until the subjet momentum balance ($\sqrt{y}$) requirement is reached:

$$\sqrt{y} \equiv \min(p_{T,j_1}, p_{T,j_2}) \frac{\Delta R(j_1,j_2)}{m_0} > 0.2,$$

(6.1)

where $p_{T,j_1}$ and $p_{T,j_2}$ are the transverse momenta of subjets $j_1$ and $j_2$ respectively, $\Delta R(j_1,j_2)$ is the distance between $j_1$ and $j_2$, and $m_0$ is the invariant mass of the pair of subjets. The selected pair of subjets is then filtered by clustering the original topological cluster constituents of the selected pair of subjets using the C/A algorithm with a small distance parameter $R_r = 0.3$. Only the three ($n_r = 3$) highest-$p_T$ subjets resulting from this reclustering are kept. (If fewer than three jets after the reclustering, then all constituents are kept.) The jet resulting from this filtering procedure is hereafter called the groomed large-$R$ jet.

6.4 Event Selection

6.4.1 Basic Selection

An event is required to have at least two reconstructed Cambridge-Aachen (C/A) jets with a radius of 1.2 and must pass the mass-drop filtering procedure. These two jets must be in the region $|\eta| < 2.0$ so that the forward calorimeter is not used and that there is good overlap with the inner tracker which is used both in the boson tagging requirements and for assessment of the systematic uncertainties. A cut is also placed on the rapidity difference between these jets, $|y_1 - y_2| < 1.2$, which offers discrimination between the more central $s$-channel signal process and the mainly $t$-channel QCD background. Also a cut on the $p_T$ asymmetry of the two leading jets, $\frac{p_{T1} - p_{T2}}{p_{T1} + p_{T2}} < 0.15$, is used to reject events where one of the jets is badly mismeasured or is not due to the primary $pp$ collision.
6.4.2 Boson-Tagged Jet

In this analysis, very loose cuts are used for the filtering stage, and no mass-drop cut is applied, as indicated in Table 6.1. When the iterative procedure stops, the remaining constituents are reclustered with the C/A algorithm using radius parameter $R_t = 0.3$, and the leading three subjets are retained ($n_r = 3$).

To improve the discrimination in this analysis, the remaining jets are also tagged with three additional boson tagging requirements. First, the jets are required to have $(\sqrt{y} \equiv \min \left(p_T(j_1), p_T(j_2)\right) \frac{\Delta R(j_1, j_2)}{m_0} \geq 0.45)$ where $\sqrt{y}$ is the value from the filtering process. This cut was chosen based on the signal and background efficiencies calculated from the first plot in Figure 6.1. Second, the number of charged-particle tracks associated with the original, ungroomed jet is required to be $n_{trk} < 30$. That is because energetic gluon jets are typically composed of more hadrons. This cut is based on the second plot in Figure 6.1, which shows the number of tracks matched to jets for selected jets in signal and background simulated events. Third, the jet is then required to be within a mass window. A narrow mass window is defined to be in the mass range 26 GeV centered at either 82.4 GeV or 92.8 GeV where the mass distributions of the $W$ and $Z$ jets peak in simulation, as shown by the distributions in the third plot in Figure 6.1. To enhance the statistical sample and simplify the treatment in some of the studies and cross checks described in the following sections, a wider mass window of $60 < m_j < 110$ GeV is also used. This is not expected to change the properties of the jets. The determination of the mass window is performed using a strategy based on an asymptotic approximation to the statistical significance, $Z_{llr}$, defined as

$$Z_{llr} = \sqrt{2((s+b)\log(1+s/b) - s)}, \quad (6.2)$$

where the signal $s$ is assumed to have a signal strength of one with respect to the nominal hypothesis cross section. More details can be found in the internal note [62].
Table 6.1: Parameters for the mass-drop filtering algorithm used to groom C/A jets.

<table>
<thead>
<tr>
<th>Filtering parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\sqrt{y}$</td>
<td>0.2</td>
</tr>
<tr>
<td>$\mu$</td>
<td>1</td>
</tr>
<tr>
<td>$R_t$</td>
<td>0.3</td>
</tr>
<tr>
<td>$n_t$</td>
<td>3</td>
</tr>
</tbody>
</table>

**Figure 6.1**: The distribution of variables used in boson tagging for signal and background simulation. The signal and background distributions are normalized to the same area.

### 6.4.3 Summary

The selection requirements are summarized in Table 6.2. Figure 6.2 shows the selection efficiencies as a function of the resonance masses for EGM $W' \rightarrow WZ$ and bulk $G_{RS} \rightarrow WW$ and $ZZ$ for simulated events with resonance mass within 10%
of the nominal signal mass. The efficiency applying only requirements for the event
topology and that after boson tagging requirement are compared as shown in the plots.

Table 6.2: Event selection requirements.

<table>
<thead>
<tr>
<th>Label</th>
<th>Meaning</th>
</tr>
</thead>
<tbody>
<tr>
<td>TRG</td>
<td>Event passes EF_j360_a10tcem trigger</td>
</tr>
<tr>
<td>DQC</td>
<td>Data quality requirement</td>
</tr>
<tr>
<td>JCL</td>
<td>Jet cleaning (bad-looser, ugly, BCH_CORR_CELL)</td>
</tr>
<tr>
<td>JET</td>
<td>Event has at least two ungroomed C/A 1.2 jets</td>
</tr>
<tr>
<td>FIL</td>
<td>Leading two jets survive filter</td>
</tr>
<tr>
<td>CEN</td>
<td>$</td>
</tr>
<tr>
<td>ETA</td>
<td>$</td>
</tr>
<tr>
<td>ASY</td>
<td>$\Delta &lt; 0.15$</td>
</tr>
<tr>
<td>TAG</td>
<td>Leading and subleading jets tagged as bosons with two mass windows</td>
</tr>
<tr>
<td>DJM</td>
<td>$m_{jj} &gt; 1.05$ TeV</td>
</tr>
</tbody>
</table>

Figure 6.2: Event selection efficiencies as a function of the resonance masses for EGM $W' \rightarrow WZ$ and bulk $G_{RS} \rightarrow WW$ and $ZZ$ for simulated events with resonance mass within 10% of the nominal signal mass. Left: before boson tagging, Right: after boson tagging.
6.5 Systematics Uncertainties

The systematic uncertainties that affect the shape of the $m_{jj}$ distribution and the normalization of the signal are summarized in this section. They are treated as nuisance parameters with various probability density functions described below. If the systematics only affect the normalization, it simply scales the distribution, entering as a scaling factor. The overall normalisation is a product of scale factors, each corresponding to an identified nuisance parameter. If the shape is affected by the systematics, it needs to be included when the signal histogram input is generated. Systematic uncertainties related to the background expectation are evaluated directly by the background estimation procedure using the fit errors as uncertainties.

Systematics due to detector effects are considered, as well as the uncertainties due to the Monte Carlo modelling. These are discussed in the following paragraphs.

Note that from now on $G(x|\bar{x},\sigma)$ denotes a Gaussian probability density function:

$$G(x|\bar{x},\sigma) = \frac{1}{\sqrt{2\pi}\sigma} e^{-\frac{(x-\bar{x})^2}{2\sigma^2}}$$  (6.3)

6.5.1 Luminosity Scale

The uncertainty on the integrated luminosity is taken to be $\pm 2.8\%$ [65]. It is calculated from the preliminary calibration of the luminosity scale derived from beam-separation scans performed in November 2012. Details about the methodology can be found in Ref. [66].

The probability density function for the luminosity is taken to be a gaussian with a standard deviation of 0.028:

$$f(S_{\text{int}}) = G(S_{\text{int}}|1,0.028)$$  (6.4)
6.5.2 Jet Energy Scale and Resolution

The jet energy scale (JES), or jet $p_T$ scale $\alpha$ is defined by

$$p_T = \alpha p_T^{MC}$$

(6.5)

where $p_T$ is the value expected in data and $p_T^{MC}$ denotes the value of $p_T$ reconstructed in simulation.

The jet energy scale is of great importance because any variance in the $p_T$ of a jet directly causes a shift in the dijet mass and thus move the position of the peak. As the background is steeply falling, a slight change in position of the excess can affect the number of expected background events and decrease the expected sensitivity to the signal by a large amount.

The systematic uncertainty is assessed by using track-to-calorimeter double ratios between the data and the MC simulation, as described in more detail in the tagging support note [67]. First the jet reconstruction and filtering algorithm are applied to massless inner-detector tracks. Then these tracks are matched to the calorimeter jets [68]. The ratio of the matched track jets $p_T$ to the calorimeter jets $p_T$ as a function of calorimeter jet’s $p_T$ is compared in simulation and data, from which the uncertainty of the jet $p_T$ scale is taken to be $1\pm0.02$. Treating this as Gaussian uncertainty, the probability density function of jet $p_T$ is

$$f(\alpha) = G(\alpha|1, 0.02)$$

(6.6)

Mismodelling of the jet energy resolution (JER) can change the reconstructed width of the diboson resonance. A nominal 5% JER is derived based on the width of the energy response for the MC signal jets after applying the tagging selection. Based on the recommendation from jet sub-structure group, a 20% uncertainty on this resolution is also implemented. The total effect of the degradation of the JER is therefore modeled by applying a Gaussian smearing to the $p_T$ of the reconstructed
jets. The nuisance parameter $\sigma_{rE}$ represents the uncertainty of JER, and has a Gaussian probability density function with mean zero and a standard deviation of $0.05 \times \sqrt{1.2^2 - 1^2}$:

$$f(\sigma_{rE}) = G(0, 0.05 \times \sqrt{1.2^2 - 1^2})$$ (6.7)

### 6.5.3 Jet Mass Scale and Resolution

The jet mass scale (JMS) $\alpha_m$ is defined by

$$m = \alpha_m m^{MC}$$ (6.8)

where $m$ is the value expected in data and $m^{MC}$ denotes the value in simulation.

As for the jet $p_T$ scale, the jet mass scale uncertainty is calculated by comparing the ratio of values for track and calorimeter jets in data and simulation. Figure 6.3 shows the double ratios as a function of the jet mass for jets passing the tagging selection requirements. The double ratio plots show that a 3% uncertainty on the JMS is needed and the probability density function used is

$$f(\alpha_m) = G(\alpha_m | 1, 0.03)$$ (6.9)

The nominal jet mass resolution (JMS) is extracted from the width of the $W/Z$ mass shape in a control sample.

A relatively pure sample of hadronic $W$ bosons can be obtained from $t\bar{t}$ samples, and systematic uncertainties on the mass scale and resolution can be derived thereafter. The filtered jet mass distributions in three bins of filtered jet $p_T$ are fit to the double crystal ball function described in Section 6.6 for both simulated events and data, as shown in Figure 6.4. The $W$ masses and widths that best fit the data distributions can be extracted for each filtered jet $p_T$ bin. To probe the sensitivity to the
Figure 6.3: Data-to-simulation track-to-calorimeter mean mass double ratio as a function of calorimeter jet mass. Left plot is for the standard simulation, and right plot represents the one where the tracking efficiency has been increased by its uncertainty.

Table 6.3: Comparison of the jet mass scale determined from alternative $t\bar{t}$ production models, to test the sensitivity of the fit to likely variations in the combinatoric background. In each fit, the crystal ball shape parameters $a_l, n_l, a_r,$ and $n_r$ are fixed.

<table>
<thead>
<tr>
<th></th>
<th>$M_W$</th>
<th>$\sigma_W$</th>
</tr>
</thead>
<tbody>
<tr>
<td>MC@NLO</td>
<td>$81.4 \pm 0.5$</td>
<td>$7.3 \pm 0.9$</td>
</tr>
<tr>
<td>POWHEG</td>
<td>$79.7 \pm 0.5$</td>
<td>$8.3 \pm 1$</td>
</tr>
<tr>
<td>MC@NLO</td>
<td>$81.5 \pm 0.8$</td>
<td>$6.0 \pm 1.9$</td>
</tr>
<tr>
<td>POWHEG</td>
<td>$80.6 \pm 1.1$</td>
<td>$11 \pm 0.1$</td>
</tr>
<tr>
<td>MC@NLO</td>
<td>$80.0 \pm 1.6$</td>
<td>$10.8 \pm 0.1$</td>
</tr>
<tr>
<td>POWHEG</td>
<td>$83.0 \pm 1.9$</td>
<td>$7.8 \pm 3.1$</td>
</tr>
</tbody>
</table>

Combinatoric backgrounds in $t\bar{t}$ events, MC@NLO is compared and the uncertainties are summarized in Table 6.3. The jet mass resolution scale determined by this fitting procedure can be sensitive to the constraints on the shape parameterization. This possibility was tested on this by performing tests of fit linearity. More details can be found in Section 9 in [67].

A nominal resolution of 7.5% is derived based on the width of the data and
MC agreement. Following the recommendation from jet sub-structure group, a 20% uncertainty on this resolution is also implemented. The total effect of the degradation of the JMR is therefore modeled by applying a Gaussian smearing to the mass of the reconstructed jets. The nuisance parameter $\sigma_{r_{\text{ms}}}$ represents the uncertainty of JMR, and is assigned a probability density function of Gaussian shape with mean zero and a standard deviation of $0.075 \times \sqrt{1.2^2 - 1^2}$ :
6.5.4 Momentum Balance Scale and Resolution

The momentum balance cut $\sqrt{y_f}$ is of great importance for boson jet tagging, thus a slight change affects the signal efficiency. The uncertainty is once again derived using the track-to-calorimeter jet double ratio, resulting a 2% shift assigned with the probability density function

$$f(\sigma_{y_f}) = G(0, 0.075 \times \sqrt{1.2^2 - 1^2})$$  \hspace{1cm} (6.10)

The uncertainty is input as shape systematics for convenience.

A nominal momentum balance resolution of 16% is extracted from the response of the momentum balance in MC for signal jets. Using the recommendation from the jet substructure group, a 20% uncertainty on this resolution is also implemented. The total effect of the degradation of the jet momentum balance is therefore modeled by applying a Gaussian smearing to each jet $\sqrt{y_f}$. The nuisance parameter $\sigma_{y_f}$ represents the uncertainty of jet momentum balance, and is assigned a probability density function of Gaussian shape with mean zero and a standard deviation of $0.16 \times \sqrt{1.2^2 - 1^2}$:

$$f(\sigma_{y_f}) = G(0, 0.16 \times \sqrt{1.2^2 - 1^2})$$  \hspace{1cm} (6.12)

6.5.5 Track Multiplicity Scale

The track multiplicity is not well modelled in the MC, so rather than evaluating the uncertainty using the calo-track double ratios, as for the other variables, a data driven technique is used. This scale factor is based on direct measurement of the efficiency of the track-multiplicity cut in data, fitting the $W$ boson mass peak in $V+$jets events. The resulting scale factor is $0.90 \pm 0.08$. Since the $n_{\text{trk}}$ cut is applied...
twice per event, once to each jet in the selection, a scale factor of 0.8 is applied per selected signal event, and assigned an uncertainty of 20%. The scale factor’s probability density function can be described by a truncated Gaussian function:

\[
G_t(\varepsilon | \varepsilon_0, \sigma_{\varepsilon}, \varepsilon_{\text{max}}) = \begin{cases} 
0 & \varepsilon < 0 \\
Ke^{-\frac{(\varepsilon - \varepsilon_0)^2}{2\sigma_{\varepsilon}^2}} & 0 \leq \varepsilon \leq \varepsilon_{\text{max}} \\
0 & \varepsilon > \varepsilon_{\text{max}} 
\end{cases}
\] (6.13)

with the following parameters:

\[
P(S_L) = G_t(S_t | 0.893, 0.08, 1.056).
\] (6.14)

### 6.5.6 Parton Shower Systematics

The uncertainty due to the parton shower and hadronization model is calculated by comparing the selection efficiencies obtained in simulation generated and showered by PYTHIA8 to the selection efficiencies obtained in simulation generated by PYTHIA8 and showered by HERWIG++. The selection efficiencies are compared after each cut is applied, up to but not including the cut on the number of ungroomed associated ghost tracks (since it is considered separately using a data driven technique). The largest difference of the selection efficiency can be considered as the systematic uncertainty due to the parton shower model. The systematic uncertainty is treated as a normalization systematic and it is modeled with Gaussian distribution:

\[
P(S_{ps}) = G(1, 0.05)
\] (6.15)

### 6.5.7 Summary

The probability density functions used for each of the sources of systematic uncertainty considered are summarized in Table 6.4.
Table 6.4: Summary of the systematic uncertainties and their probability distribution functions.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>probability density function</th>
<th>Meaning</th>
</tr>
</thead>
<tbody>
<tr>
<td>$S_L$</td>
<td>$G(S_L</td>
<td>1, 0.028)$</td>
</tr>
<tr>
<td>$\alpha$</td>
<td>$G(\alpha,</td>
<td>, 1, 0.02)$</td>
</tr>
<tr>
<td>$\sigma_E$</td>
<td>$G(\sigma_E</td>
<td>0, 0.05 \times \sqrt{1.2^2 - 1^2})$</td>
</tr>
<tr>
<td>$\alpha_m$</td>
<td>$G(\alpha_m</td>
<td>1, 0.03)$</td>
</tr>
<tr>
<td>$\sigma_m$</td>
<td>$G(\sigma_m</td>
<td>0, 0.075 \times \sqrt{1.2^2 - 1^2})$</td>
</tr>
<tr>
<td>$\alpha_y$</td>
<td>$G(\alpha_y</td>
<td>1, 0.02)$</td>
</tr>
<tr>
<td>$\sigma_y$</td>
<td>$G(\sigma_y,</td>
<td>0, 0.16 \times \sqrt{1.2^2 - 1^2})$</td>
</tr>
<tr>
<td>$S_t$</td>
<td>$G(S_t</td>
<td>0.89, 0.095, 1.07)$</td>
</tr>
<tr>
<td>$S_{ps}$</td>
<td>$G(S_{ps}</td>
<td>1.0, 0.05)$</td>
</tr>
</tbody>
</table>

6.6 Background Estimation

The dominant background components are dijet events arising from QCD processes. The dijet mass spectrum is expected to be smoothly falling but leading order or NLO MC predictions for QCD background are not expected to describe the data very accurately. As a result, the search strategy for diboson resonances is to look for bumps on a smoothly falling background in the dijet mass spectrum, which is based on a dijet mass function that has been shown to model dijet mass well in many previous analyses.

The dijet function is as follows:

$$\frac{dn}{dx} = p_1 (1 - x)^{p_2 + \xi p_3} x^{p_3}$$  \hspace{1cm} (6.16)

where $x = m_{jj}/\sqrt{s}$, $\sqrt{s}$ has dimension of mass and is fixed to 8 TeV. In such a way, the function is valid in the range $0 < m_{jj} < \sqrt{s}$ and goes to zero at the kinematic limit due to the collision energy. $p_1$ is a normalization factor, $p_2$ and $p_3$ are dimensionless shape parameters, and $\xi$ is a constant chosen from MC to minimize the correlation between $p_2$ and $p_3$ in the fit. The reduction of the correlation leads to the advantage
that fewer MC events are required to gain the same numerical precision, without affecting the quality of the fit.

To test the compatibility of the dijet mass function with the QCD background, we apply it to some simulated background as well as some sidebands in the data. The results of the fits are listed in Table 6.5. It shows the fits to data where one jet mass falls in the low-mass sideband ($40 < m_{j1} < 60$ GeV) and the other falls in a high-mass sideband from $110 < m_{j2} < 140$ GeV, and where both jet masses fall in the high-mass sideband. One-sided $\chi^2$ probabilities are displayed. The probabilities are calibrated using pseudo-experiments to avoid biases due to empty bins for the three data sideband fits.

The results of the fitting in the dijet mass spectrum are shown in Figure 6.5, where the fitting method is applied to data before tagging and after $WZ$, $WW$, and $ZZ$ selection. The significance shown in the inset for each bin is the difference between the data and the fit in units of the uncertainty on this difference. The spectra are compared to the signals expected for an EGM $W'$ with $m_{W'} = 2.0$ TeV or to an RS graviton with $m_{G^*} = 2.0$ TeV.

It can be seen that in these plots the dijet function describes very well all three selections except for a few bins around 2 TeV where more events are observed than predicted. In Section 6.8 I describe detailed tests of this excess to ensure that it is not caused by the detector or some other effects. There is approximately 20% overlap of the events between the three selections in the signal regions.

6.7 Limits

Upper limits on the production cross section times branching ratio of the heavy resonance are evaluated using the same statistical procedure described in detail in Chapter 8. The resulting limit plots are shown in Figure 6.6 using EGM $W'$ as benchmark model for the $WZ$ selection and using the bulk $G_{RS}$ as benchmark model
for the $WW$ and $ZZ$ selection. EGM $W' \rightarrow WZ$ for masses between 1.3 and 1.5 TeV are excluded at 95% CL while no mass region is excluded for bulk $G_{RS}$ with the sensitivity of this measurement. All of the three observed limits fall outside of the $2\sigma$ sensitivity band around 2 TeV, with the local $p_0$ values shown in Figure 6.5. The most significant excess with respect to the background expectation is seen in the $WZ$ selection with a local significance of $3.4\sigma$. The global significance is also
Table 6.5: Goodness-of-fit for maximum-likelihood fits of the background model to the dijet mass distribution in simulated events.

<table>
<thead>
<tr>
<th>Sample</th>
<th>$\chi^2$/nDOF</th>
<th>Probability</th>
</tr>
</thead>
<tbody>
<tr>
<td>PYTHIA dijet events</td>
<td>24.6/22</td>
<td>0.31</td>
</tr>
<tr>
<td>dijet events</td>
<td>15.9/22</td>
<td>0.82</td>
</tr>
<tr>
<td>Data with $110 &lt; m_{j1} \leq 140$ GeV and $40 &lt; m_{j2} \leq 60$ GeV</td>
<td>12.1/11</td>
<td>0.79</td>
</tr>
<tr>
<td>Data with $40 &lt; m_j \leq 60$ GeV for both jets</td>
<td>19.8/13</td>
<td>0.56</td>
</tr>
<tr>
<td>Data with $110 &lt; m_j \leq 140$ GeV for both jets</td>
<td>5.0/6</td>
<td>0.91</td>
</tr>
</tbody>
</table>

computed by taking into account all the mass ranges with all the three channels, which corresponds to $2.5 \sigma$.

Figure 6.6: Upper limits on the cross section times branching ratio as a function of $m_{W'}$ for the WZ selection (a), of $m_{G_{RS}}$ for the ZZ selection (b), of $m_{G_{RS}}$ for the ZZ selection (c) at 95% C.L., and local $p_0$ for three selections (d).
6.8 Cross Checks

The following section describes extensive cross checks to determine if the excess is caused by some unaccounted for detector effect.

6.8.1 Pileup Dependence Studies

Proton-proton collisions in addition to the collision of interest are referred to as “pileup”, which presents a serious challenge to physics analyses at the LHC. The main effect of pileup on jet is the increase of the noise term of the jet energy resolution, causing local fluctuations. In-time pileup refers to the additional $p$-$p$ collisions occurring in the same bunch-crossing as the collision of interest; out-of-time pileup is the additional $p$-$p$ collisions occurring in bunch-crossings just before and after the collision of interest. These extra collisions can have an effect when detectors are sensitive to several bunch-crossings or their electronics integrate over more than 25 ns.

Boson tagging and $m_{jj}$ pileup dependence studies are summarized in this section.

6.8.1.1 Pileup Effect on Boson Tagging

The boson tagging selection efficiency as a function of the mean number of the interactions per crossing $\langle \mu \rangle$ (out-of-time pileup) and the number of vertices $N_{PV}$ (in-time pileup) is studied in this section. In figure 6.7, 6.8, 6.9, and 6.10, five dijet mass ($p_T^j$) bins are used: 1.0-1.2 TeV, 1.2-1.4 TeV, 1.4-1.6 TeV, 1.6-1.8 TeV and above 1.8 TeV. The lower mass region, corresponding to the red mark, shows a clear positive correlation with the out of time pileup while the other mass regions are quite flat. Focusing on the selection cut efficiency for these five dijet mass bins, we apply a series of selection cuts—the baseline cut, the cut on $n_{trk}$, the wide mass window cut on first jet, the wide mass window cut on second jet and the subjet momentum balance cut $y_f$, as shown in Figure 6.11 and 6.12. It is obvious that there is visible
pileup dependence for the cut on jet mass. But after the posterior $y_f$ cut (subjet momentum balance cut) is applied, the slope disappears.

Figure 6.13 and 6.14 show the number of clusters in the leading jet in the five dijet mass bins before the $y_f$ cut and after the $y_f$ cut separately. Before the subjet momentum balance cut, the shapes corresponding to different $\langle \mu \rangle$ bins peak at different places, showing pileup dependence. However after the cut, the distributions shift and become quite similar.

These studies all point to the conclusion that our choice of the selection cuts for the analysis control the pileup effect on the boson tagging in a reasonable way. More discussion of the effect of pileup on the rest of the analysis is presented in the next subsection.

6.8.1.2 $m_{jj}$ Pileup Dependence

Figure 6.15 and 6.16 shows the $m_{jj}$ spectrum in three $\langle \mu \rangle$ and $N_{PV}$ bins for data. These two distributions test the effects of total pileup and in-time pileup respectively.
Figure 6.8: Boson tagging selection cuts (wide mass window cut on first jet) efficiency as a function of the mean number of the interactions per crossing.

Figure 6.9: Boson tagging selection cuts (wide mass window cut on second jet) efficiency as a function of the mean number of the interactions per crossing.
Figure 6.10: Boson tagging selection cuts (subjct moment balance cut $y_f$) efficiency as a function of the mean number of the interactions per crossing.

A Kolmogorov test (KS test) is run to assess the compatibility in shape between the three distributions which corresponding to three different levels of pileup activity. It is known that the value of the returned value for a KS test for binned data will be shifted slightly higher than the probability from a comparison of the unbinned data. However, as long as the binned width is small compared to any significant physical effect (for example, the experimental resolution), the binning will not be a huge issue. The result for the different $\langle \mu \rangle$ binned distributions and $N_{PV}$ binned distributions can be seen in Table 6.6 and Table 6.7. The returned function value is the probability of test, and much less than 1 means not compatible. Therefore, the result of the test shows that the three shapes are in reasonable agreement.

Proton bunches injected into the beam are ordered in bunch trains (groups of bunches) followed by some empty bunches. The effect of out-of-time pileup was also tested by looking at the $m_{jj}$ compatibility as a function of the distance from the bunch train front (measured in filled bunches). Figure 6.17 shows the $m_{jj}$ distribu-
Figure 6.11: Selection cut efficiency as a function of the mean number of the interactions per bunch crossing in five dijet mass bins (top left: 1.0-1.2 TeV, top right: 1.2-1.4 TeV, middle left: 1.4-1.6 TeV, middle right: 1.6-1.8 TeV, bottom left: above 1.8 TeV) for data. The cuts are applied in sequence as follows: baseline selection, number of tracks associated to jets with ghost association procedure, wide mass window cut on first jet, wide mass window cut on second jet, subjet moment balance cut $y_f$.

6.8.2 Jet Event Displays

To get a clearer view of the substructure and pileup conditions of the selected jets, we make jet display plots using the JetAnalysisDisplay tool. Different event objects, such as the jet and its constituents (either clusters or truth MC particles), tracks, and subjets, are displayed by the tool, providing a versatile platform for looking at jet substructure. It is extremely useful for showing how the pileup effects the jets.
Figure 6.12: Selection cut efficiency as a function of the number of vertices in five dijet mass bins (top left: 1.0-1.2 TeV, top right: 1.2-1.4 TeV, middle left: 1.4-1.6 TeV, middle right: 1.6-1.8 TeV, bottom left: above 1.8 TeV) for data. The cuts are applied in sequence as follows: baseline selection, number of tracks associated to jets with ghost association procedure, wide mass window cut on first jet, wide mass window cut on second jet, subjet moment balance cut $y_f$.

after the tagging procedure.

Two stages of jet filtering are introduced: the simple mass drop (SMD) stage and the reclustered mass drop (RMD) stage. As discussed in Section 6.4.1 and 6.4.2, the mass drop filtering technique is applied in the boson identification process. The SMD jet is the jet obtained at the end of the mass-drop procedure. Then we filter the SMD jet by reclustering the remaining constituents to obtain the RMD jet. The leading $n_r$ jets are retained ($n_r = 3$).

We select several events around the excess and consider them in more detail by looking at the jet display plots. In Figure 6.18 and Figure 6.19 active jet area is plotted for event 75855145 at the SMD stage and RMD stage respectively. The
Figure 6.13: Number of clusters in jet in 5 dijet mass bins at different pileup level before \( y_f \) cut.

Figure 6.14: Number of clusters in jet in 5 dijet mass bins at different pileup level after \( y_f \) cut.

black block represents the leading large-R jet and the magenta block the subleading one. Among them, the grey solid blocks are the subjets that are obtained after the mass drop-filtering step. Clusters are shown color coded according to their \( p_T \), tracks are shown as black dots and pileup tracks are shown in azure crosses. Comparing to Figure 6.18, Figure 6.19 clearly shows the merged subjets after the
Figure 6.15: $m_{jj}$ spectrum in three $\langle \mu \rangle$ bins for data.

Figure 6.16: $m_{jj}$ spectrum in three $N_{PV}$ bins for data.
Table 6.6: KS test result for spectrum in three $\langle \mu \rangle$ bins for data.

<table>
<thead>
<tr>
<th>$\mu$ range</th>
<th>KS result</th>
</tr>
</thead>
<tbody>
<tr>
<td>[0, 18] and [18, 24]</td>
<td>0.630836</td>
</tr>
<tr>
<td>[0, 18] and [24, 50]</td>
<td>0.520089</td>
</tr>
<tr>
<td>[18, 24] and [24, 50]</td>
<td>0.99285</td>
</tr>
</tbody>
</table>

Table 6.7: KS test result for spectrum in three $N_{PV}$ bins for data.

<table>
<thead>
<tr>
<th>$N_{PV}$ range</th>
<th>KS result</th>
</tr>
</thead>
<tbody>
<tr>
<td>[0, 11] and [11, 15]</td>
<td>0.293046</td>
</tr>
<tr>
<td>[0, 11] and [15, 50]</td>
<td>0.223008</td>
</tr>
<tr>
<td>[11, 15] and [15, 50]</td>
<td>1</td>
</tr>
</tbody>
</table>

filtering procedure. From all these plots, it can been seen that pileup is not very severe in the subjet region, which demonstrates that it would not affect the excess too much. More displays for other events are shown in Figure 6.20 and Figure 6.21.
Figure 6.17: $m_{jj}$ spectrum in three distance from train front bins for data.

Table 6.8: KS test result for spectrum in three distance from train front bins for data.

<table>
<thead>
<tr>
<th>$\mu$ range</th>
<th>KS result</th>
</tr>
</thead>
<tbody>
<tr>
<td>[0, 11] and [12, 24]</td>
<td>0.961</td>
</tr>
<tr>
<td>[0, 11] and [25, 38]</td>
<td>0.686</td>
</tr>
<tr>
<td>[12, 24] and [25, 38]</td>
<td>0.788</td>
</tr>
</tbody>
</table>
Figure 6.18: Jet displays for event 75855145 at SMD stage.
Figure 6.19: Jet displays for event 75855145 at RMD stage.

Figure 6.20: Jet displays for event 60808462.
Figure 6.21: Jet displays for event 114169287.
Combined Search for Diboson Resonances

7.1 Introduction

To get better sensitivity to new diboson resonances, a combination of four statistically independent searches for diboson resonances is conducted and interpreted as constraints on three benchmark models, the EGM, the bulk-RS model, and the HVT model. In addition to the fully hadronic search described in the previous chapter, diboson resonances decaying to $WZ$, $WW$ and $ZZ$ in the $\ell\nu\ell\ell$ ($l = \mu, e$), $\ell\ell qq$, $\ell\nu qq$ final states are discussed. The EGM interpretation combine the results of all four channels described above. For the bulk RS interpretation, only the $\ell\ell qq$, $\ell\nu qq$, and $JJ$ channels are combined since only a neutral resonance is modeled in this benchmark scenario.

Both ATLAS and CMS have set limits on new diboson resonances with $pp$ collision data delivered by the LHC. CMS has set limits on the production of resonances using a combination of $\ell\ell qq$, $\ell\nu qq$ and fully hadronic final states [69]. This study shown here improves the cross section limits at low mass, and extends the mass exclusion for both $W'$ and bulk $G^*$ with respect to previous searches.
Figure 7.1: Acceptance times efficiency for the different analyses entering into the EGM $W'$ combination.

7.2 Search Channels

Four separate channels enter the combination that exploit different decays of the $W$ and/or $Z$ bosons: $WZ \rightarrow \ell\nu\ell\ell$ (full-leptonic), $WV \rightarrow \ell\nu qq$ (semi-leptonic), $ZV \rightarrow \ell\ell qq$ (semi-leptonic), and $VV \rightarrow qqqq$ (full-hadronic).

The four analyses are mutually exclusive and therefore the channels are statistically independent. They are made orthogonal by exclusive requirements on the number of leptons. The $\ell\ell qq$ analysis only uses events with two leptons with $p_T > 20$ GeV, while the $\ell\nu qq$ analysis vetos events with more than one lepton, and the all-hadronic analysis vetoes events with any leptons. Details about the three leptonic analyses, which were not part of the searches carried out in this thesis, are provided below.
7.2.1 $WZ \rightarrow \ell \nu \ell \ell$

This search is sensitive to charged resonances decaying to $WZ$ with the $\ell \nu \ell \ell$ [70] final state. For this analysis [71], events are selected requiring three electrons or muons. Two of them must have opposite sign and are the decaying product of a $Z$ boson. The $z$ component of the neutrino momentum is inferred by constraining the $W$ mass, which leads to a full reconstruction of the $WZ$ four-momentum. Two signal regions—low-mass and high-mass regions, are defined. In the low-mass region the lepton not associated with the $Z$-boson decay and the $E_T^{\text{miss}}$ have to meet $\Delta \phi(\ell, E_T^{\text{miss}}) > 1.5$, however the cut is reversed in the high-mass region. The high-mass region is used to test all signal masses under study, while the low-mass region is only used to constrain...
W' masses less than 400 GeV.

The invariant mass of the $WZ$ pair is the final discriminant. The leading background is SM $WZ$ production, which is modeled with MC simulated samples. An extrapolation is performed to high $WZ$ masses to compensate for poor MC statistics for $m(WZ) > 1$ TeV. The $WZ$ background is reduced by angular requirements on the separation of the $W$ and $Z$. The other backgrounds are from $ZZ$ and other events with three real leptons, which are also modeled with simulation. $Z$+jets background and other processes producing fewer than three prompt leptons are estimated in situ by rescaling events containing leptons that pass relaxed identification criteria.

Plots of the observed data compared to the expected background are shown in Fig. 7.3 for the low-mass signal region and in Fig. 7.4 for the high-mass signal region.

7.2.2 $WV \rightarrow \ell \nu qq$

This search is sensitive to charged resonances decaying to $WZ$ and neutral resonances decaying to $WW$ with the $\ell \nu qq$ [72] final state. For this analysis [73], events are selected requiring a single lepton (electron or muon), and either two small-$R$ jets or a single large-$R$ jet. The mass of the small-$R$ jets pair or the single merged jet is required to be consistent with the $W$ and $Z$ boson mass. Those events with $b$-tagged jet are discarded to suppress backgrounds from $t\bar{t}$ and single-top production. Two exclusive channels are defined, one electron channel and one muon channel, and three orthogonal signal regions are constructed for each channel: the high-$p_T$ resolved region (events with two small-$R$ jets that satisfy higher $p_T$ thresholds on the dijet); the low-$p_T$ resolved region (events with two small-$R$ jets that do not satisfy the high-$p_T$ resolved selection) and the boosted region (events with a large-$R$ jet).

The mass of the $WV$ distribution is studied as the discriminant. The dominant background is $W$+jets, which are reweighted using $p_T(\ell\nu)$ measured in the sideband region to improve the data/MC agreement. The shape of the $m(WV)$ distribution is
obtained from simulation, and the normalization is derived by scaling MC to the data in the sidebands of the dijet (if two small-$R$ jets) or large-$R$ jet mass distribution. Other backgrounds, for instance, $Z$+jets, $t\bar{t}$, and other processes producing real leptons are also estimated with simulated samples; the contributions from multijet events with fake or non-prompt leptons are estimated *in situ* by rescaling the events containing leptons that pass relaxed identification criteria.
Figure 7.4: Results of the $W' \rightarrow WZ \rightarrow \ell\nu\ell\ell$ search in the high-mass signal region in each of the four flavor channels [71].

Plots of the expected backgrounds and observed data are shown in Fig. 7.5 for the low-$p_T$ resolved signal region, in Fig. 7.6 for the high-$p_T$ resolved signal region, and in Fig. 7.7 for the boosted signal region.
Figure 7.5: Results of the $WV \to \ell qq$ search in the low-$p_T$ resolved signal region for the electron (a) and muon (b) channels [73].

Figure 7.6: Results of the $WV \to \ell qq$ search in the high-$p_T$ resolved signal region for the electron (a) and muon (b) channels [73].
Figure 7.7: Results of the $WV \rightarrow \ell\nu qq$ search in the merged signal region for the electron (a) and muon (b) channels [73].

7.2.3 $ZV \rightarrow \ell\ell qq$

This search is sensitive to charged resonances decaying to $WZ$ and neutral resonances decaying to $ZZ$ with the $\ell\ell qq$ [74] final state. For this analysis [75], events are selected requiring two electrons or muons, and either two small-$R$ jets or a single large-$R$ jet. The mass of the small-$R$ jets pair or the single merged jet is required to be consistent with the $W$ and $Z$ boson mass. Three signal regions are defined, based on the $p_T$ of the leptons and the jets: the high-$p_T$ resolved region (events with two small-$R$ jets that satisfy higher $p_T$ thresholds on the dijet), the low-$p_T$ resolved region (events with two small-$R$ jets that do not satisfy the high-$p_T$ resolved selection) and the merged region (events with a large-$R$ jet).

The three regions are orthogonal. Event selection for the merged region has the highest priority, followed by the high-$p_T$ resolved region. Table 7.1 gives the $p_T$ thresholds defined for these regions.

The final discriminant is the invariant mass of the $ZV$ system. The dominant
Table 7.1: Definitions of the three signal regions defined in the $\ell\ellqq$ search.

<table>
<thead>
<tr>
<th>Signal region</th>
<th>$p_T^{\ell\ell}$</th>
<th>$p_T^{qq}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Merged</td>
<td>$&gt; 400$ GeV</td>
<td>$&gt; 400$ GeV</td>
</tr>
<tr>
<td>High-$p_T$ resolved</td>
<td>$&gt; 250$ GeV</td>
<td>$&gt; 250$ GeV</td>
</tr>
<tr>
<td>Low-$p_T$ resolved</td>
<td>$&gt; 100$ GeV</td>
<td>$&gt; 100$ GeV</td>
</tr>
</tbody>
</table>

Figure 7.8: Results of the $ZV \rightarrow \ell\ellqq$ search in the low-$p_T$ resolved signal region for the $ee$ (a) and $\mu\mu$ (b) channels [75].

Background is $Z$+jets, which is modeled with simulation. The normalization of the simulated $Z$+jets samples is scaled to the data in the sidebands of the dijet or large-$R$ jet mass distribution. A bin-by-bin shape correction for the $Z$+jets template is derived using the same region. Subdominant backgrounds from dibosons, $t\bar{t}$, and other processes producing two real leptons are also estimated with simulated samples. Backgrounds from processes such as $W$+jets or QCD multijets that do not produce two prompt leptons are negligible.

Plots of the expected backgrounds and observed data are shown in Fig. 7.8 for the low-$p_T$ resolved signal region, in Fig. 7.9 for the high-$p_T$ resolved signal region, and in Fig. 7.10 for the merged signal region.
Figure 7.9: Results of the $ZV \rightarrow \ell\ell qq$ search in the high-$p_T$ resolved signal region for the $ee$ (a) and $\mu\mu$ (b) channels [75].

Figure 7.10: Results of the $ZV \rightarrow \ell\ell qq$ search in the merged signal region for the $ee$ (a) and $\mu\mu$ (b) channels [75].
7.2.4 Summary

Figure 7.1 summarizes the acceptance times efficiency for a $W'$ in the different analyses as a function of the $W'$ mass. Figure 7.2 summarizes the acceptance times efficiency for a $G^*$ in the different analyses as a function of the $G^*$ mass, considering only $G^* \rightarrow VV$ decays.

7.3 Data Samples

The data analyzed in this study used the dataset collected in 2012 by ATLAS detector at $\sqrt{s} = 8$ TeV. The triggers used are a single-lepton (electron or muon) trigger with a $p_T$ threshold of 24 GeV, and a single large-radius jet trigger with a $p_T$ threshold of 360 GeV. The total integrated luminosity is $\int L \, dt = 20.3 \, fb^{-1}$. The uncertainty on the integrated luminosity is 2.8%.

7.4 Monte Carlo Samples

7.4.1 Signal

The signal samples are generated with pole masses between 200 and 2500 GeV in 100 GeV steps. The $W'$ signal events are generated with Pythia8 [76], setting the coupling scale factor $c = 1$; while the bulk $G^*$ signal events are produced by CalcHEP [77] with $k/\bar{M}_{Pl} = 1.0$. The factorization and renormalization scales are modeled with Pythia8 in both cases and are set to the generated resonance mass. The CTEQ6L1 [63] (MSTW2008LO [78]) parton distribution functions (PDFs) are used for the $G^*$ ($W'$) signal. The widths, leading order cross sections and branching ratios for the $W'$ and bulk $G^*$ signal samples for a few mass points are listed in Table 7.2.
Table 7.2: Leading order cross-sections, branching ratios and widths for $W'$ and bulk $G^*$ signals at $\sqrt{s} = 8$ TeV for selected mass points.

<table>
<thead>
<tr>
<th>m [TeV]</th>
<th>$\sigma(W')$ [fb]</th>
<th>BR($W' \rightarrow WZ$) [%]</th>
<th>$\Gamma_{W'}$ [GeV]</th>
<th>$\sigma(G^*)$ [fb]</th>
<th>BR($G^* \rightarrow WW$) [%]</th>
<th>BR($G^* \rightarrow ZZ$) [%]</th>
<th>$\Gamma_{G^*_{\text{RS}}}$ [GeV]</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.5</td>
<td>3198.1</td>
<td>1.6</td>
<td>18.0</td>
<td>3106.0</td>
<td>34</td>
<td>18</td>
<td>18.4</td>
</tr>
<tr>
<td>1.0</td>
<td>152.2</td>
<td>1.3</td>
<td>36.0</td>
<td>56.0</td>
<td>19</td>
<td>10</td>
<td>55.4</td>
</tr>
<tr>
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<td>18.7</td>
<td>1.3</td>
<td>54.0</td>
<td>3.14</td>
<td>17</td>
<td>8</td>
<td>89.5</td>
</tr>
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<td>3.14</td>
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<td>73.3</td>
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<td>16</td>
<td>8</td>
<td>122.5</td>
</tr>
<tr>
<td>2.5</td>
<td>0.69</td>
<td>1.2</td>
<td>90.7</td>
<td>0.032</td>
<td>16</td>
<td>8</td>
<td>155.0</td>
</tr>
</tbody>
</table>

7.4.2 Backgrounds

Unlike the fully-hadronic search described in Chapter 7, the leptonic channel searches rely heavily on MC modeling of their dominant backgrounds. The $W+$jets and $Z+$jets backgrounds are modeled by the SHERPA generator [79] with CT10 PDFs [80], PYTHIA6 with CTEQ6L1 PDFs, and ALPGEN [81] to estimate the systematic uncertainty. The $W+$jets and $Z+$jets production cross sections are scaled to NNLO calculations [82].

The $t$-channel single-top events are generated by AcerMC [83] with CTEQ6L1 PDFs and PYTHIA6 for hadronisation. The $t\bar{t}$ production, $s$-channel single-top and $Wt$ processes are modeled with the NNLO [84, 85] generator. CT10 PDFs are used. Hadronisation is performed by HERWIG [86] and the underlying event is modeled by JIMMY [87]. The $t\bar{t}$ production cross section is calculated at NNLO in QCD including resummation of next-to-next-to-leading logarithmic soft gluon terms with Top++2.0 [88, 89, 90, 91, 92, 93].

QCD multi-jet events are simulated with PYTHIA6, HERWIG, and POWHEG [94, 95] interfaced to PYTHIA6.

The SM diboson production is modeled with the HERWIG generator and CTEQ6L1 PDFs. An exception is the $\ell\nu\ell'\ell'$ channel, which uses POWHEG interfaced to PYTHIA6. The diboson production cross sections are normalized to predictions at next-to-
leading-order accuracy [96]. Additional diboson samples for the $\ell\nu qq$ channel are produced with Sherpa.

Generated events are processed with the ATLAS detector simulation [97] based on the GEANT4 package [98]. Pileup events are simulated with Pythia8.

7.5 Object Reconstruction and Event Selection

The search channels participating in the combined search share some common reconstructed objects. Specifically, reconstructed jets, electrons, muons, and missing transverse momentum are discussed below.

7.5.1 Jets

Jets are reconstructed by combining topological clusters in the calorimeter system [58], which are calibrated in energy using the local calibration weighting scheme [61] and are considered massless.

7.5.1.1 Small-$R$ jets

Hadronically decaying vector bosons with low-$p_T$ ($\lesssim 450$ GeV) are reconstructed with two jets. Each jet is formed with the anti-$k_T$ algorithm [16] with a distance parameter $R = 0.4$, which is referred to as a small-$R$ jet. Both the $\ell\nu qq$ and $\ell\ell qq$ channels use the small-$R$ jets, among which only jets within $|\eta| < 2.8$ (2.1) and $p_T > 30$ GeV are considered. The scalar summed $p_T$ of the tracks associated to the primary vertex must account for at least 50% of the scalar summed $p_T$ of all tracks associated to the jet for those jets with $p_T < 50$ GeV, to reject jets dominated by pileup activity. Jets containing hadrons from $b$-quarks are identified using a multivariate $b$-tagging algorithm [99].
7.5.1.2 Large-$R$ jets

Hadronically decaying vector bosons with high-$p_T$ ($\gtrsim 400$ GeV) are reconstructed with a single merged jet due to the collimated nature of their decay products. The large-$R$ jets, denoted with $J$, are first constructed using the Cambridge–Aachen (C/A) algorithm [100] with a distance parameter $R = 1.2$ and then groomed according to the procedure described in Chapter 6.

7.5.1.3 Boson jet identification

As in the fully-hadronic search, additional cuts are imposed on large-$R$ jets to identify jets from hadronically decaying bosons. The $\sqrt{y}$ cuts are the same in all analyses. In the one lepton and dilepton channels, mass windows of $65 < m_j < 105$ GeV and $70 < m_j < 110$ GeV, where $m_j$ represents the jet mass, are used respectively. In the fully hadronic channel, mass windows of $69.4 < m_j < 95.4$ GeV and $79.8 < m_j < 105.8$ GeV are used to select $W$ or $Z$ boson respectively. Finally, the fully-hadronic channel cuts on the number of tracks ghost-associated to the large-$R$ jet ($n_{trk} < 30$).

7.5.2 Leptons

An introduction to electron and muon reconstruction at ATLAS can be found in Section 5.4 and 5.5. These analyses use the high-$p_T$ electron reconstruction (seeded from the calorimeter) in the central region, excluding the transition region between the barrel and the endcap calorimeters ($1.37 < |\eta| < 1.52$). The transverse energy $E_T$ of an electron is defined using the energy of the EM cluster and the pseudorapidity of the track since the angular resolution of the inner detector is better than the calorimeters, but the calorimeter has superior energy resolution, especially at high electron transverse momenta. Only electrons with $E_T > 25$ GeV are considered. Electrons are required to pass the tight identification criteria for reconstruction of $W \rightarrow e\nu$ candidates, while $Z \rightarrow ee$ are reconstructed with electrons that satisfy...
the medium identification criteria, as described in Chapter 5. Muon candidates are required to be combined, reconstructed within the range $|\eta| < 2.5$ [101]. Only muons with $p_T > 25$ GeV are considered.

Backgrounds due to misidentified leptons and non-prompt leptons are suppressed by lepton isolation requirements based on tracking and calorimeter information. However, in some cases, the isolation requirement needs to be modified. For example, when a $Z$ boson is produced in the boosted regime, the leptons from its decay can be within the isolation angle. Therefore, to maintain a high efficiency for diboson resonances with masses above 1 TeV, contributions from the additional lepton are excluded when computing the isolation discriminants for dilepton objects of the same flavor with high-$p_T$ ($\gtrsim 600$ GeV). Details of the lepton isolation criteria including the parameters of the isolation discriminants and the maximum thresholds applied, can be found in Table 7.3.

<table>
<thead>
<tr>
<th>Channel</th>
<th>Calorimeter</th>
<th>ID</th>
<th>Cone Size $\Delta R$</th>
</tr>
</thead>
<tbody>
<tr>
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<td>$&lt;0.16$</td>
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<td>$&lt;0.15$</td>
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<tr>
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<td>$&lt;0.14$</td>
<td>$&lt;0.15$</td>
<td>0.2</td>
</tr>
</tbody>
</table>

7.5.3 Missing $E_T$

The missing transverse energy in this analysis is used to estimate the transverse momentum of the neutrino in relevant sub-channels. The magnitude of the missing transverse momentum $E_T^{\text{miss}}$ is calculated as the magnitude of the vectorial sum of the transverse momenta of the electrons, muons, jets as well as calibrated energy deposits in the calorimeter which are not associated to a high-$p_T$ jet. [102].
7.6 Systematic Uncertainties

The sources of systematic uncertainties are object reconstruction, the energy scale and resolution of leptons, missing $E_T$, small-$R$ and large-$R$ jets. Other uncertainties related to luminosity and event generation are also taken into account. Details about the uncertainties considered and their effect on the estimated signal and background yields are documented in corresponding publications [103], [104], [105] and [106]. Although the selections from different search channels coming into the combination are statistically independent, some common systematics on the objects used, the integrated luminosity, and the signal and background simulation, across different search channels, introduce correlated effects in the signal and background expectations. Therefore, it is important to properly treat these correlations in the combination analysis. Table 7.4 lists all systematic uncertainties that are considered to be fully correlated across different channels. Note that this table is not a complete list of all systematics considered in the likelihood. Only those nuisance parameters common to more than one decay channel are shown here.

7.7 Statistical Treatment and Combination

The statistical treatment for the analysis as well as the analysis described in Chapter 7, is based on a frequentist significance test using a likelihood ratio as a test statistic [107]. A binned maximum-likelihood (ML) fit to the diboson mass spectrum in each channel is performed using histogram templates derived from MC with data driven estimates. The signal strength $\mu$, which is the scale factor on the cross section based on the signal hypothesis, is one of the parameters of interest. $\mu=0$ corresponds to the background-only hypothesis, and $\mu=1$ corresponds to the nominal signal hypothesis. The effect of systematic uncertainties are modeled with nuisance parameters $\theta_k$ which are unknown a priori. The corresponding probability density
Table 7.4: List of systematic uncertainties considered to be correlated across different channels. The × symbol indicates that the systematic is included in the likelihood for the given channel. The † marker indicates that the nuisance parameter is largely constrained via control regions or sidebands within an individual channel, and is thus not highly correlated between channels; for instance, the luminosity uncertainty for ℓνqq does not affect the W+jets normalization, as this is determined in situ using data.

<table>
<thead>
<tr>
<th>Uncertainty name</th>
<th>JJ</th>
<th>ℓℓjj</th>
<th>ℓνjj</th>
<th>ℓνℓℓ</th>
<th>Comments</th>
</tr>
</thead>
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<tr>
<td>ISRFSR</td>
<td>×</td>
<td>×</td>
<td>×</td>
<td></td>
<td>ISR/FSR for gluons and quarks</td>
</tr>
<tr>
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<td>×†</td>
<td>×†</td>
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<td></td>
<td>Luminosity</td>
</tr>
<tr>
<td>PDF</td>
<td>×</td>
<td>×</td>
<td>×</td>
<td>×</td>
<td>PDF</td>
</tr>
<tr>
<td>ElecLER</td>
<td>×</td>
<td>×</td>
<td>×</td>
<td></td>
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</tr>
<tr>
<td>ElecLESLowPt</td>
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<td></td>
<td></td>
<td></td>
<td>Electron trigger efficiency</td>
</tr>
<tr>
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<td></td>
<td></td>
<td>Muon ID track scale</td>
</tr>
<tr>
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<td></td>
<td></td>
<td></td>
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<tr>
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<td>×</td>
<td></td>
<td></td>
<td></td>
<td>Muon reco efficiency</td>
</tr>
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<td></td>
<td></td>
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</tr>
<tr>
<td>MuonID</td>
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<td></td>
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<td>Muon ID track scale</td>
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<td>MuonLES</td>
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<td></td>
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<td></td>
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</tr>
<tr>
<td>MuonMS</td>
<td>×</td>
<td></td>
<td></td>
<td></td>
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<tr>
<td>MuonReco</td>
<td>×</td>
<td></td>
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<td></td>
<td>Muon reco efficiency</td>
</tr>
<tr>
<td>MuonTrig</td>
<td>×</td>
<td></td>
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<td></td>
<td>Muon trigger efficiency</td>
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<td>METSoftResol</td>
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<td></td>
<td>JER for large-R jets</td>
</tr>
<tr>
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<td></td>
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<td>JER for large-R jets</td>
</tr>
<tr>
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<td></td>
<td></td>
<td>JMR for large-R jets</td>
</tr>
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<td></td>
<td></td>
<td></td>
<td>JMS for large-R jets</td>
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<td>YR for large-R jets</td>
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<td></td>
<td></td>
<td>YS for large-R jets</td>
</tr>
<tr>
<td>JER</td>
<td>×</td>
<td></td>
<td></td>
<td></td>
<td>Jet energy resolution</td>
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<td>×</td>
<td></td>
<td></td>
<td></td>
<td>Jet energy scale parameter</td>
</tr>
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<td></td>
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<td></td>
<td>JES MC closure test uncertainty</td>
</tr>
<tr>
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<td></td>
<td></td>
<td></td>
<td>JES η-intercalibration parameter</td>
</tr>
<tr>
<td>JSEEtaInter2</td>
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<td></td>
<td></td>
<td>JES η-intercalibration parameter</td>
</tr>
<tr>
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<td></td>
<td></td>
<td></td>
<td>JES flavor dependence</td>
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<tr>
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<td></td>
<td></td>
<td>JES flavor response</td>
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<tr>
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<td></td>
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<td></td>
<td></td>
<td></td>
<td>JES pileup parameter</td>
</tr>
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<td></td>
<td></td>
<td>JES pileup parameter</td>
</tr>
<tr>
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<td></td>
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<td>JES pileup parameter</td>
</tr>
<tr>
<td>JESSinglePart</td>
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<td></td>
<td></td>
<td></td>
<td>JES pileup parameter</td>
</tr>
</tbody>
</table>
functions $f_k(\theta_k)$ help to constrain $\theta_k$. The likelihood function is the product of the Poisson probabilities for all bins in all channels:

$$L = \prod_i \prod_j \text{Pois}(n_{ij}^{\text{obs}} | n_{ij}^{\text{exp}}(\mu, \theta_k)) \prod_k f_k(\theta_k)$$  \hspace{1cm} (7.1)$$

where $i$ represents the sub-channels that go into the combination, and $j$ stands for the bin in the invariant mass spectrum.

To get an estimate on the value of $\mu$, the likelihood ratio is then computed:

$$\lambda(\mu) = \frac{L(\mu, \hat{\theta})}{L(\hat{\mu}, \hat{\theta})},$$  \hspace{1cm} (7.2)$$

where $\mu$ is the signal strength, and $\hat{\theta}$ is a vector of the nuisance parameter values that maximize the likelihood $L$ for the specified $\mu$ (conditional ML) and thus is a function of $\mu$, $\hat{\mu}$ and $\hat{\theta}$ are the signal strength and nuisance parameter values that maximize $L$ (unconditional ML). The presence of nuisance parameters broadens the distribution of the profile likelihood as a function of $\mu$, and thus leads to loss of information about $\mu$.

From Eq. (7.2), one can see that $0 \leq \lambda \leq 1$, with $\lambda$ close to 1 implying good agreement between data and the hypothesized value of $\mu$. Therefore, it is convenient to use a test statistic defined as:

$$t_\mu = -2 \ln \lambda(\mu)$$  \hspace{1cm} (7.3)$$

with a high value of $t_\mu$ implying incompatibility between data and $\mu$.

To extract limits we performed two binned ML fits, one in which the signal is allowed to float, and one in which the signal is fixed to zero. The likelihood ratio

$$\tilde{\lambda}(\mu) = \begin{cases} 
\frac{L(\mu, \hat{\theta})}{L(0, \hat{\theta})}, & \text{if } \hat{\mu} < 0, \\
\frac{L(\mu, \hat{\theta})}{L(\hat{\mu}, \hat{\theta})}, & \text{if } \hat{\mu} \geq 0, 
\end{cases}$$  \hspace{1cm} (7.4)$$

was used to calculate the test statistic $\tilde{q}_\mu$ is given by

$$\tilde{q}_\mu = \begin{cases} 
-2 \ln \bar{\lambda}(\mu) & \text{if } \hat{\mu} \leq \mu, \\
0 & \text{if } \hat{\mu} > \mu.
\end{cases} \quad (7.5)$$

In the case that a significant positive signal is observed, the statistic $\tilde{q}_\mu$ described above is used to test $\mu = 0$ in a class of model where we assume $\mu \geq 0$. Rejecting $\mu = 0$ hypothesis leads to the discovery of a new signal:

$$q_0 = \begin{cases} 
-2 \ln \lambda(0) & \text{if } \hat{\mu} \geq 0, \\
0 & \text{if } \hat{\mu} < 0.
\end{cases} \quad (7.6)$$

The exclusion limits at the 95% confidence level is extracted using CLs prescription [108], which is a statistical method for setting upper limits and an approximate confidence in the signal-only hypothesis. It is a frequentist method and the limits are defined by means of error probabilities. Upper limits derived with the CLs method always contain the zero value of the parameter and hence the coverage probability at this point is always 100%. The level of agreement between observed data and the given hypothesis $H$ is quantified by the $p$ value, defined as the probabilility of the background only prediction to produce a signal like deviation at least as large as that observed in the data. With $t_\mu$ known, $p$ value is then calculated as:

$$p_\mu = \int_{t_{\mu,\text{obs}}}^{\infty} f(t_\mu|\mu) \, dt_\mu, \quad (7.7)$$

In order to make the computation more efficient and less time-consuming, some of the systematics are considered to be negligible for the combination. Two tests are performed to decide if a systematic uncertainty is significant: normalization and shape tests. First, the normalization of the nominal spectra is compared with the systematically shifted spectra. If the deviation is greater than 0.5 $\sigma$ of the statistical error, this systematic uncertainty is considered. Second, a KS test is performed to
compare the shapes of the nominal distribution and the systematically shifted one. If the returned value of the KS test is less than 0.1, that systematic uncertainty is included as a shape uncertainty.

The combination is performed in the mass range 0.2 to 3.0 TeV. The profile likelihood is performed simultaneously over multiple selection regions for different channels with the correlations for nuisance parameters taken into account. For different mass points, the selection regions that are sensitive are different. For the case of $W'$, all four sub-channels are combined; while for the case of $G^*$, only $\ell\ell qq$, $\ell\nu qq$ and the fully hadronic channels contribute to the combined search. The sensitive regions and channels contributed are listed in Table 7.5 and Table 7.6 for $W'$ and $G^*$ samples respectively.

7.8 Results

The limit from the full combination is shown in Fig. 7.11 for the EGM $W'$ signal model and in Fig. 7.13 for the bulk $G^*$ model. A comparison of the median expected limits for all decay channels in the EGM $W'$ and the bulk $G^*$ signal models is shown in Fig. 7.12 and Fig. 7.14 respectively. The observed $p_0$ value is shown in Fig. 7.15 for the combination and in Fig 7.16 for the individual channels. The combined results show that the most significant deviation is observed in the EGM $W'$ model at 2 TeV with a local significance of 2.5 standard deviations. In the sub-channels, the most significant deviation from the background expectation is found in the fully hadronic channel at the same mass in the $W'$ channel with a local significance of 3.4 standard deviations and a global significance of 2.5 standard deviations. For the other sub-channels, no significant deviations from the background expectation are observed.

The limits are obtained using leading order cross sections and the asymptotic
Table 7.5: Selection regions used for each $W'$ mass point for limit setting.

<table>
<thead>
<tr>
<th>$M[TeV]$</th>
<th>$qqqq$</th>
<th>$\ell\nuqq$</th>
<th>$\ell\ellqq$</th>
<th>$\ell\ell\ell$</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.2</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>low-mass,high-mass</td>
</tr>
<tr>
<td>0.3</td>
<td>-</td>
<td>low-$p_T$</td>
<td>low-$p_T$</td>
<td>low-mass,high-mass</td>
</tr>
<tr>
<td>0.4</td>
<td>-</td>
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<td>low-$p_T$</td>
<td>low-mass,high-mass</td>
</tr>
<tr>
<td>0.5</td>
<td>-</td>
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<td>low-$p_T$</td>
<td>low-mass,high-mass</td>
</tr>
<tr>
<td>0.6</td>
<td>-</td>
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<td>low-$p_T$, high-$p_T$</td>
<td>low-mass,high-mass</td>
</tr>
<tr>
<td>0.7</td>
<td>-</td>
<td>low-$p_T$, high-$p_T$</td>
<td>low-$p_T$, high-$p_T$</td>
<td>low-mass,high-mass</td>
</tr>
<tr>
<td>0.8</td>
<td>-</td>
<td>low-$p_T$, high-$p_T$, boosted</td>
<td>low-$p_T$, high-$p_T$</td>
<td>low-mass,high-mass</td>
</tr>
<tr>
<td>0.9</td>
<td>-</td>
<td>high-$p_T$, boosted</td>
<td>low-$p_T$, high-$p_T$, boosted</td>
<td>low-mass,high-mass</td>
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<td>1.0</td>
<td>-</td>
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<td>high-$p_T$, boosted</td>
<td>low-mass,high-mass</td>
</tr>
<tr>
<td>1.1</td>
<td>boosted</td>
<td>boosted</td>
<td>high-$p_T$, boosted</td>
<td>low-mass,high-mass</td>
</tr>
<tr>
<td>1.2</td>
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<td>boosted</td>
<td>high-$p_T$, boosted</td>
<td>low-mass,high-mass</td>
</tr>
<tr>
<td>1.3</td>
<td>boosted</td>
<td>boosted</td>
<td>high-$p_T$, boosted</td>
<td>low-mass,high-mass</td>
</tr>
<tr>
<td>1.4</td>
<td>boosted</td>
<td>boosted</td>
<td>high-$p_T$, boosted</td>
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</tr>
<tr>
<td>1.5</td>
<td>boosted</td>
<td>boosted</td>
<td>high-$p_T$, boosted</td>
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</tr>
<tr>
<td>1.6</td>
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<td>boosted</td>
<td>high-$p_T$, boosted</td>
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</tr>
<tr>
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<td>boosted</td>
<td>boosted</td>
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</tr>
<tr>
<td>1.8</td>
<td>boosted</td>
<td>boosted</td>
<td>high-$p_T$, boosted</td>
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</tr>
<tr>
<td>1.9</td>
<td>boosted</td>
<td>boosted</td>
<td>high-$p_T$, boosted</td>
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</tr>
<tr>
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<td>high-mass</td>
</tr>
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<td>high-$p_T$, boosted</td>
<td>high-mass</td>
</tr>
<tr>
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<td>high-mass</td>
</tr>
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</tr>
<tr>
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<td>high-$p_T$, boosted</td>
<td>high-mass</td>
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<td>-</td>
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<td>-</td>
<td>-</td>
<td>-</td>
</tr>
</tbody>
</table>

approximation is used\(^1\). For EGM $W''$, the resulting lower limit after the full combination is observed to be 1.81 TeV, with an expected limit of 1.81 TeV. The most stringent limit of an individual analysis comes from the $\ell\nuqq$ channel, which is 1.59 TeV.

For the bulk $G^*$ the observed combined limit is 810 GeV, with an expected limit of 790 GeV. The most stringent individual limit is from the $\ell\ellqq$ channel with 740 GeV.

\(^1\) The asymptotic approximation is expected to be valid up to masses of around 2 TeV. In the limit of high statistics, the profile likelihood ratio is a parabolic function of the parameter with a minimum at $\hat{\mu}$. In the high-mass region there are insufficient data statistics to use this approximation. The result for high mass region is checked by running toys.
Figure 7.11: The combined limit on the EGM $W'$ using all channels. Results include statistical and systematic uncertainties.

7.9 HVT Interpretation

The cross section for $W'$ production in the HVT model is very similar to that in the EGM model, and the production mechanism is also identical. Therefore, the constraints on the HVT $W'$ can be obtained from the EGM $W'$ results by scaling the EGM cross section to the HVT cross section at different points in the HVT parameter space, assuming the same acceptance and efficiency. The results are presented as constraints on HVT parameters $\{(g^2/g_V)c_F, g_Vc_H\}$ in the range where the extrapolation from the the EGM $W'$ results is expected to be valid. The constraints on the HVT $W'$ signal are shown in Fig. 7.17 for the individual analyses, and in Fig 7.18 for the leptonic and full combinations. Since the HVT model predicts both a $W'$ and a $Z'$ with very similar masses, the constraints shown here are weaker than they should be when considering the full model.
7.10 HVT vs Graviton

A study using MC generator output without detector effects has been conducted, comparing the kinematics between HVT and graviton samples (HVT neutral $V \rightarrow WW \rightarrow qqqq$ Vs. Graviton $\rightarrow WW \rightarrow qqqq$). The motivation for the study is the similarity observed in the decay modes for the two models. Certain weight factors, shown in Table 7.7 have been applied to the HVT samples to offset the higher $p_T$ of the graviton samples on the whole. The event selection efficiencies are compared between these two models at each selection stage. Table 7.8 presents the numbers after each cut—the fiducial cut, the dijet balance cut, the momentum balance cut and the jet mass cut. Slight differences can be seen from these numbers, indicating possible explanations due to the production mechanism of these two different samples. It can be understood qualitatively by considering that the $G^*$ and HVT are produced...
Figure 7.13: The combined limit on the graviton using all channels. Results include statistical and systematic uncertainties.

predominantly by gg and qqbar processes, respectively, and hence more radiation in gg→G* events could cause a shift in the jet mass, resulting in a lower efficiency for the G*. This was indeed seen in the truth comparison between the G* and W' signals in Figure 7.19, 7.20 and 7.21 which show jet mass, subjet momentum balance and jet p_T for leading jet and subleading jet respectively. The G* distributions have a longer tail towards high mass, which appears to be consistent with this.
Figure 7.14: A comparison of the median expected limits for all decay channels in the graviton signal model.
Table 7.6: Selection regions used for each $G^*$ mass point for limit setting.

<table>
<thead>
<tr>
<th>$M$ [TeV]</th>
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<th>$\ell\ell qq$</th>
</tr>
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</tr>
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</tr>
<tr>
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</tr>
<tr>
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<tr>
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Table 7.7: Weight factor applied to the HVT samples.

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<th>GeV</th>
<th>HVT(1.2)</th>
<th>G(1.2)</th>
<th>HVT(1.8)</th>
<th>G(1.8)</th>
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<th>G(2.4)</th>
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<td>boson Pt(Aver)</td>
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<td>609</td>
<td>819</td>
<td>880</td>
<td>1078</td>
<td>1115</td>
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<td>weighted factor</td>
<td>1.107</td>
<td>1.074</td>
<td>1.034</td>
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Figure 7.15: The observed $p_0$-values for the combination of the EGM $W'$ and graviton signal hypotheses. Results include statistical and systematic uncertainties.

Table 7.8: Event selection efficiencies applying selection cuts sequentially.

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<th>%</th>
<th>HVT(1.2)</th>
<th>G(1.2)</th>
<th>HVT(1.8)</th>
<th>G(1.8)</th>
<th>HVT(2.4)</th>
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<td>$</td>
<td>\Delta y</td>
<td>&lt;1.2$</td>
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<td>82.80</td>
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<td>Dijet balance</td>
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<td>Momentum balance</td>
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<td>62.58</td>
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<td>RMD jet mass</td>
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<td>51.96</td>
<td>55.69</td>
<td>50.10</td>
<td>56.00</td>
<td>49.33</td>
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</table>
Figure 7.16: The observed $p_0$-values for the sub-channels of the EGM $W'$ and graviton signal hypotheses: $JJ$ (top left), $\ell\nuqq$ (top right), $\ell\ellqq$ (bottom left) and $\ell\nu\ell\ell$ (bottom right). Results include statistical and systematic uncertainties.
Figure 7.17: Observed 95% CL exclusion contours in the HVT parameter space \{\((g_v^2/g_F)_{cF}, g_V c_H}\) for mass points 1.0 TeV, 1.5 TeV, 1.8 TeV and 2.0 TeV for each sub-channel. The benchmark model parameters are also shown: \(A(g_V = 1), A(g_V = 3)\) and \(B(g_V = 3)\).
**Figure 7.18:** Observed 95% CL exclusion contours in the HVT parameter space \{(g^*/g_V)c_F, g_Vc_H\} for mass points 1.0 TeV, 1.5 TeV, 1.8 TeV and 2.0 TeV for combination of channels with leptons and combination of all channels. The benchmark model parameters are also shown: \(A(g_V = 1), A(g_V = 3)\) and \(B(g_V = 3)\).

**Figure 7.19:** Kinematics comparison between the \(G^*\) and \(W'\) signals for pole mass 1.2 GeV.
Figure 7.20: Kinematics comparison between the $G^*$ and $W'$ signals for pole mass 1.8 GeV.

Figure 7.21: Kinematics comparison between the $G^*$ and $W'$ signals for pole mass 2.4 GeV.
Conclusions

The LHC provides a unique yet challenging environment for conducting a search for new heavy resonances at the TeV scale. As the luminosity, and center of mass collision energy rise, it gets difficult to measure the hadronic decays of objects. However, it is extremely useful to exploit these channels since hadronic decays contribute to a large fraction of boson decays. This thesis has presented the work for conducting the search for heavy resonance using full hadronic channel as well as the combined results.

The work of Chapter 6 is published in JHEP with the title of *Search for high-mass diboson resonances with boson-tagged jets in proton-proton collisions at $\sqrt{s} = 8$ TeV with the ATLAS detector* [103]. The paper has been cited by more than 160 papers in only a few months, and several new theoretical models have been inspired by the result, such as an anomaly free leptophobic left-right model [109], leptophobic $U(1)_{LP}$ models from string theories [110], and extra $Z'$s and $W'$s in heterotic-string derived models [111].

The work of Chapter 7 is accepted by PLB with the title of *Combination of searches for WW, WZ, and ZZ resonances in pp collisions at $\sqrt{s} = 8$ TeV with the*
Figure 8.1: The limits set on the bulk RS graviton through (a) the $ZZ$ and (b) the $WW$ channels, and on the HVT through (c) the $Z' \rightarrow WW$ and (d) the $W' \rightarrow WZ$ channels [112].
ATLAS detector [113].

The results of an updated search for a heavy resonance decaying to dibosons in the all hadronic channel using 3.2 fb\(^{-1}\) of new 13 TeV \(pp\) collision data collected by the ATLAS experiment at the LHC in 2015 [112] are presented. The data are in good agreement with the background expectations at all masses. Figure 8.1 shows that upper limits are set in the mass region 1.2 to 2.2 TeV on the cross-section for production of the HVT model through the \(W' \rightarrow WZ\) and \(Z' \rightarrow WW\) modes, and the bulk GRS model through the \(ZZ\) and \(WW\) decays. The HVT model is excluded in the mass region 1.38 to 1.6 TeV at 95\% CL from the \(W' \rightarrow WZ\) search. The most extreme local \(p_0\)-value is in the HVT \(W' \rightarrow WW\) channel around 2 TeV. The results show great compatibility with the Run 1 results.

It is important to note that this analysis work represents contributions from everyone in the analysis team and is not done all by myself, as this is the nature of big collaborations. The work on ATLAS Run 1 Pythia8 tunes in Chapter 4, the in-situ validation of the jet mass scale using top events, jet event displays, and cross checks on the pileup dependence of the boson tagging rates in Chapter 6, the HVT and graviton samples’ comparison and the HVT interpretation in Chapter 7 was carried out solely by me. In Chapter 7 I carried out the statistical treatment and combination limit calculation. The plots that I did not produce have corresponding references.
Bibliography


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Biography

Lei Li was born 26 April 1988 in China. She grew up in Qingdao, Shandong Province, eastern China and received her bachelor’s degree in physics from Shandong University in 2010. Then she went to Durham, North Carolina, USA and completed her Ph.D. studies in experimental high energy physics at Duke University in 2016.