SEARCH FOR NEW NEUTRAL HIGH-MASS RESONANCES DECAYING INTO MUON PAIRS WITH THE ATLAS DETECTOR

by

Simon Viel

B.Sc., Université Laval, 2008

A THESIS SUBMITTED IN PARTIAL FULFILLMENT OF THE REQUIREMENTS FOR THE DEGREE OF

Doctor of Philosophy

in

THE FACULTY OF GRADUATE AND POSTDOCTORAL STUDIES
(Physics)

THE UNIVERSITY OF BRITISH COLUMBIA
(Vancouver)

August 2014

© Simon Viel, 2014
Abstract

The question of physics beyond the Standard Model remains as crucial as it was before the discovery of a Higgs boson at the Large Hadron Collider, as the theoretical and experimental shortcomings of the Standard Model remain unresolved. Indeed, theoretical problems such as the hierarchy of energy scales, the Higgs mass fine-tuning and the large number of postulated parameters need to be addressed, while the experimental observations of dark matter, dark energy and neutrino masses are not explained by the Standard Model. Many hypotheses addressing these issues predict the existence of new neutral high-mass resonances decaying into muon pairs.

This dissertation documents a search for this process using 25.5 fb$^{-1}$ of proton-proton collision data collected by the ATLAS experiment in Run-I of the Large Hadron Collider. After evaluating the performance of the detector for reconstructing muons at very high momentum, the event yields observed as a function of the invariant mass of muon pairs are compared with expected values from Standard Model processes.

The observed yields are found to be in good agreement with Standard Model predictions, and no significant excess of events is found. New gauge bosons with couplings to fermions equal to these of the Standard Model Z boson and with masses lower than 2.53 TeV are therefore excluded at 95% confidence level. A statistical combination with the results of the search for the same particle decaying into electron pairs yields a lower mass limit of 2.90 TeV at 95% confidence level. Limits are also placed in the context of two classes of models inspired by Grand Unification Theories: gauge theories with the $E_6$ symmetry group, as well as Minimal $Z'$ Models.
Preface

The research presented in this dissertation is based on the experimental data of the ATLAS experiment at the Large Hadron Collider. Thousands of researchers based at more than 170 institutions from 38 countries participate in the ATLAS collaboration.

All of the dissertation text was written by me, and not taken directly from previously published sources. Only versions of the text in Sections 7.2, 7.3, A.3.6 and A.3.7 have previously appeared in internal ATLAS documentation. All figures and tables for which a reference or figure credit is not indicated in the caption represent my own work, often in collaboration with other researchers named below. Figures labeled “ATLAS”, “ATLAS Preliminary” and event displays with the ATLAS logo are public material released by the ATLAS collaboration. Figures labeled “ATLAS work in progress” are previously unpublished material derived from ATLAS real and/or simulated data.

Chapters 2, 3 and 4 present introductory material, with all sources referenced in the text. I participated in commissioning the Transition Radiation Tracker described in Section 4.3.3 and contributed to developing and maintaining the data quality monitoring system described in Section 4.8.3. Section 5.1 also presents introductory material from the indicated references. From the material in Section 5.2 I have collaborated with L. Chevalier and H. Wang to produce Figure 5.3, and I have produced Figure 5.6.

I was one of the lead investigators for the analysis presented in Section 5.3 and Chapters 6, 7, and 8. This work led directly to four publications, first with the full dataset collected by ATLAS in 2010 [Phys. Lett., B700:163, 2011], then with the initial part of the dataset collected in 2011 [Phys. Rev. Lett., 107:272002, 2011], the full dataset collected in 2011 [JHEP, 1211:138, 2012] and finally with the full dataset collected in 2012 [Accepted by Phys. Rev. D, 2014].

In addition, many other published analyses from the ATLAS collaboration have also made use of the primary muon selection described in Section 5.3 and in particular four of these have also used the methods described in Chapters 6 and 7 searching for a non-resonant excess of events using the same datasets collected in 2010 [Phys. Rev., D84:011101, 2011], in 2011 [Phys. Lett., B712:40, 2012; Phys. Rev., D87:015010, 2013] and in 2012 [Submitted to Eur. Phys. J. C, 2014].

My main roles toward these publications were as follows:

• For the first two publications above, I was the main data analyst in the dimuon channel.
• I then served as co-leader of the analysis with the full 2011 dataset, coordinating a team of 70 researchers from 27 institutes. I was primarily responsible from the dimuon channel, while S. Heim was co-leader for the dielectron channel. In addition to refining the techniques used in the analysis, we have set stringent limits on a wide variety of theoretical models.

• For the analysis with the full 2012 dataset, I continued to make crucial collaborations to the analysis, indicated below.

• I contributed importantly to writing these four publications, along with internal supporting documentation, and addressed comments during the review process.

• I also contributed to estimating the background and signal predictions as well as the related uncertainties for the four other publications listed.

In detail, my original contributions to the material presented in the main body of this dissertation are as follows:

• I contributed to developing the muon selection in Section 5.3.1 and characterizing the performance of the detector for these muons, as member of a task force coordinated by S. Willocq.

• I carried out the studies described in Section 5.3.2 for the analysis with the full 2011 dataset, starting from code by D. Fortin. I collaborated with J. Coggeshall to replicate these studies for the analysis with the 2012 dataset.

• I contributed to developing the event selection described in Section 6.1 and implemented it, starting from code by I. Nugent. The yield histograms in this section were produced in collaboration with E. Laisné, and the event display was produced in collaboration with L. Chevalier.

• I performed the background and signal estimates described in Section 6.2, with the following inputs provided by collaborators:
  
  – Real and simulated data prepared centrally by members of the ATLAS collaboration;
  – Theoretical corrections to the \(Z/\gamma^*\) differential cross section calculated by T. Nunnemann, J. Kretzschmar and U. Klein;
  – Fits to the tails of the sub-leading backgrounds performed by S. Heim and E. Fitzgerald;
  – Cosmic ray background estimate designed by P. Wagner;
  – Signal template re-weighting method designed by A. Kotwal and O. Stelzer-Chilton, with interference effects included by W. Fedorko, S. Heim, N. Hod and A. Kotwal.

• I performed the comparison of data with background expectations in Section 6.3 in the first three rounds of analysis, starting from code by D. Hayden. Variants of my code were used for the analysis with the 2012 dataset by R. Daya and E. Fitzgerald.
• I developed the techniques used to evaluate theoretical and experimental systematic uncertainties discussed in Chapter 7, using inputs by T. Nunnemann and U. Klein for the theoretical uncertainties on the $Z/\gamma^*$ yields. The experimental uncertainties were evaluated in collaboration with D. Fortin and O. Stelzer-Chilton for the analyses with the 2010 and 2011 datasets, and with J. Coggeshall and E. Fitzgerald for the analysis with the 2012 dataset.

• I developed and managed the statistical framework used in Chapter 8, starting from code by B. Stelzer interfacing the Bayesian Analysis Tools. W. Fedorko and S. Heim also contributed to this framework.


• The expected limits presented in Chapter 9 are derived from work carried out primarily by T. Hryn’ova, U. Klein, J. Kretzschmar and me.

In Appendix A, Sections A.1 and A.2 present introductory material relevant to the work presented in Section A.3. My original contributions to this analysis in development are the following:

• I investigated possible gains in sensitivity with the Matrix Element method, starting from code mainly by D. Schouten, and also by M. Bluteau, P. Chang, B. Stelzer and K. van Nieuwkoop.

• I optimized the event pre-selection and the Boosted Decision Tree input variables, in collaboration with B. Cerio, K. McLean, D. Schouten, B. Stelzer and K. van Nieuwkoop.

• I performed modelling studies of variables and correlations, and estimated the $Z/\gamma^*$ background in both search channels, starting from code by D. Schouten.

• I quantified the effect of systematic uncertainties on the shape of the discriminant variable, starting from code by M. Venturi and T. Lenz.

• I contributed to the statistical framework of the analysis group, and used it to quantify the performance of the multivariate analysis developed in comparison to previous results.
# Table of Contents

Abstract ................................................................. ii

Preface ................................................................. iii

Table of Contents ...................................................... vi

List of Tables ........................................................... x

List of Figures .......................................................... xii

Glossary ................................................................. xix

Acknowledgements ...................................................... xxii

1 Introduction .......................................................... 1

2 The Standard Model of Particle Physics ......................... 3
   2.1 The Standard Model Lagrangian ................................ 3
      2.1.1 Forces and Matter .................................... 3
      2.1.2 The Englert-Brout-Higgs-Guralnik-Hagen-Kibble Mechanism .... 6
      2.1.3 Electroweak Interactions Revisited .................... 8
      2.1.4 Summary .............................................. 10
   2.2 Successes of the Standard Model ............................. 10
      2.2.1 Asymptotic Freedom and Infrared Slavery .......... 11
      2.2.2 Calculation of Cross Sections and Decay Rates .... 13
   2.3 Limitations of the Standard Model ......................... 22
      2.3.1 Quantum Gravity .................................... 22
      2.3.2 Dark Matter and Dark Energy ....................... 22
      2.3.3 Charge-Parity Violation and Matter-Antimatter Asymmetry .... 24
      2.3.4 Neutrino Masses and Flavour Oscillation .......... 24
<table>
<thead>
<tr>
<th>Section</th>
<th>Title</th>
<th>Location</th>
</tr>
</thead>
<tbody>
<tr>
<td>5</td>
<td>Muons at Very High Momentum in ATLAS</td>
<td>63</td>
</tr>
<tr>
<td>5.1</td>
<td>Muon Reconstruction in ATLAS</td>
<td>64</td>
</tr>
<tr>
<td>5.2</td>
<td>High-Momentum Muon Performance</td>
<td>65</td>
</tr>
<tr>
<td>5.3</td>
<td>Dedicated Very-High Momentum Muon Selection</td>
<td>74</td>
</tr>
<tr>
<td>5.3.1</td>
<td>Selection of 3-Station Muons</td>
<td>74</td>
</tr>
<tr>
<td>5.3.2</td>
<td>Selection of 2-Station Muons</td>
<td>76</td>
</tr>
<tr>
<td>6</td>
<td>Event Selection and Comparison of Data with Standard Model Expectations</td>
<td>83</td>
</tr>
<tr>
<td>6.1</td>
<td>Event Selection</td>
<td>83</td>
</tr>
<tr>
<td>6.2</td>
<td>Background and Signal Expectation</td>
<td>89</td>
</tr>
<tr>
<td>6.2.1</td>
<td>Simulated Samples</td>
<td>89</td>
</tr>
<tr>
<td>6.2.2</td>
<td>Data-Driven Background Estimates</td>
<td>96</td>
</tr>
<tr>
<td>6.2.3</td>
<td>Generation of Signal Templates</td>
<td>100</td>
</tr>
<tr>
<td>6.3</td>
<td>Comparison of Data with Background Expectations</td>
<td>104</td>
</tr>
<tr>
<td>6.3.1</td>
<td>Kinematics of the Dimuon System</td>
<td>105</td>
</tr>
<tr>
<td>6.3.2</td>
<td>Kinematics of Individual Muons</td>
<td>105</td>
</tr>
<tr>
<td>6.3.3</td>
<td>Missing Transverse Energy Distributions</td>
<td>117</td>
</tr>
<tr>
<td>7</td>
<td>Systematic Uncertainties</td>
<td>121</td>
</tr>
<tr>
<td>7.1</td>
<td>Overview</td>
<td>121</td>
</tr>
<tr>
<td>7.1.1</td>
<td>Theoretical Uncertainties</td>
<td>122</td>
</tr>
<tr>
<td>7.1.2</td>
<td>Experimental Uncertainties</td>
<td>123</td>
</tr>
<tr>
<td>7.2</td>
<td>Parton Distribution Function and QCD Uncertainties on Signal Cross Sections</td>
<td>126</td>
</tr>
<tr>
<td>7.3</td>
<td>Parton Distribution Function and QCD Uncertainties on the $Z/\gamma^*$ Cross Section</td>
<td>128</td>
</tr>
<tr>
<td>7.3.1</td>
<td>Parton Distribution Function Variations</td>
<td>129</td>
</tr>
<tr>
<td>7.3.2</td>
<td>Parton Distribution Function Set Choice, QCD Scale and $\alpha_S$</td>
<td>133</td>
</tr>
<tr>
<td>8</td>
<td>Statistical Methods and Results</td>
<td>137</td>
</tr>
<tr>
<td>8.1</td>
<td>Signal Search</td>
<td>137</td>
</tr>
<tr>
<td>8.1.1</td>
<td>Local $p$-Values</td>
<td>138</td>
</tr>
<tr>
<td>8.1.2</td>
<td>Global $p$-Values</td>
<td>142</td>
</tr>
<tr>
<td>8.2</td>
<td>Limits on $Z'_{SSM}$ and $E_6 Z'$ Bosons</td>
<td>146</td>
</tr>
<tr>
<td>8.3</td>
<td>Limits on Minimal $Z'$ Models</td>
<td>150</td>
</tr>
<tr>
<td>9</td>
<td>Outlook</td>
<td>155</td>
</tr>
<tr>
<td>10</td>
<td>Conclusion</td>
<td>157</td>
</tr>
<tr>
<td></td>
<td>Bibliography</td>
<td>159</td>
</tr>
</tbody>
</table>
A Search for Vector Boson Fusion $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ ................................. 172
A.1 Status of Higgs Boson Observations at the Large Hadron Collider ................. 172
A.2 Analysis Goals and Strategy ............................................................................. 175
A.2.1 Physical Objects .......................................................................................... 177
A.2.2 Backgrounds ............................................................................................... 178
A.2.3 Topology of VBF $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ Events ......................... 179
A.2.4 Boosted Decision Trees .............................................................................. 182
A.2.5 Matrix Element Method ............................................................................. 184
A.3 Contributions to the Analysis ....................................................................... 185
A.3.1 Statistical Framework .............................................................................. 185
A.3.2 Event Pre-Selection Optimization ............................................................... 186
A.3.3 Boosted Decision Tree Optimization ........................................................... 187
A.3.4 Modelling Studies ...................................................................................... 188
A.3.5 Correlation Studies .................................................................................... 194
A.3.6 Background Estimation and Systematic Uncertainties for $Z/\gamma^* \rightarrow \tau^+\tau^-$ . 195
A.3.7 Background Estimation and Systematic Uncertainties for $Z/\gamma^* \rightarrow e^+e^-$, $\mu^+\mu^-$ in the Same-Flavour Channel ................................................................. 196
A.3.8 Effect of Object Systematic Uncertainties on the BDT Score Shape .......... 199
A.3.9 Comparison of Analysis Techniques .......................................................... 200
A.3.10 Matrix Element Method Investigation ..................................................... 206
A.4 Conclusion and Outlook ............................................................................... 209
# List of Tables

<table>
<thead>
<tr>
<th>Table Number</th>
<th>Description</th>
<th>Page</th>
</tr>
</thead>
<tbody>
<tr>
<td>2.1</td>
<td>Gauge bosons in the Standard Model</td>
<td></td>
</tr>
<tr>
<td>2.2</td>
<td>Fermions in the Standard Model, in flavour basis</td>
<td></td>
</tr>
<tr>
<td>3.1</td>
<td>Values of $\gamma'$ and $\theta_{\text{Min}}$ for three specific models: $Z'_{B-L}$, $Z'_X$ and $Z'_R$</td>
<td></td>
</tr>
<tr>
<td>3.2</td>
<td>Previous experimental limits at 95% CL on the mass of new gauge bosons $Z'$</td>
<td></td>
</tr>
<tr>
<td>5.1</td>
<td>Muon $p_T$ resolution parameters for simulated samples generated in 2011 at $\sqrt{s} = 7\text{TeV}$</td>
<td></td>
</tr>
<tr>
<td>5.2</td>
<td>Muon $p_T$ resolution parameters for simulated samples generated in 2012 at $\sqrt{s} = 8\text{TeV}$</td>
<td></td>
</tr>
<tr>
<td>5.3</td>
<td>Muon momentum smearing constants used in analyses at $\sqrt{s} = 7\text{TeV}$. The effect of $S^\text{ID}_1$ is neglected.</td>
<td></td>
</tr>
<tr>
<td>5.4</td>
<td>Muon momentum smearing constants used in analyses at $\sqrt{s} = 8\text{TeV}$.</td>
<td></td>
</tr>
<tr>
<td>5.5</td>
<td>$p^{\text{MS}^<em>}_2$ parameters measured in data and simulation at $\sqrt{s} = 7\text{TeV}$ and corresponding smearing parameters $S^{\text{MS}^</em>}_2 = P^{\text{MS}^<em>}<em>2,</em>{\text{data}} \odot P^{\text{MS}^</em>}<em>2,</em>{\text{MC}}$.</td>
<td></td>
</tr>
<tr>
<td>5.6</td>
<td>$p^{\text{MS}^*}_2$ parameters measured in data and simulation at $\sqrt{s} = 8\text{TeV}$.</td>
<td></td>
</tr>
<tr>
<td>6.1</td>
<td>Summary of the simulated samples for the analysis using the full dataset at $\sqrt{s} = 7\text{TeV}$</td>
<td></td>
</tr>
<tr>
<td>6.2</td>
<td>Summary of the simulated samples for the analysis using the full dataset at $\sqrt{s} = 8\text{TeV}$</td>
<td></td>
</tr>
<tr>
<td>6.3</td>
<td>Expected and observed number of events for the analysis at $\sqrt{s} = 7\text{TeV}$</td>
<td></td>
</tr>
<tr>
<td>6.4</td>
<td>Expected and observed number of events for the analysis at $\sqrt{s} = 8\text{TeV}$</td>
<td></td>
</tr>
<tr>
<td>7.1</td>
<td>Summary of systematic uncertainties on the expected number of events for the search using data collected in 2011 at $\sqrt{s} = 7\text{TeV}$.</td>
<td></td>
</tr>
<tr>
<td>7.2</td>
<td>Summary of systematic uncertainties on the expected numbers of events for the search using data collected in 2012 at $\sqrt{s} = 8\text{TeV}$.</td>
<td></td>
</tr>
<tr>
<td>7.3</td>
<td>Uncertainty on $Z'$ cross sections due to PDF variations at 90% CL at $\sqrt{s} = 7\text{TeV}$.</td>
<td></td>
</tr>
</tbody>
</table>
## List of Figures

<table>
<thead>
<tr>
<th>Figure</th>
<th>Description</th>
<th>Page</th>
</tr>
</thead>
<tbody>
<tr>
<td>2.1</td>
<td>Particle content of the Standard Model</td>
<td>10</td>
</tr>
<tr>
<td>2.2</td>
<td>Gauge couplings of the Standard Model as a function of energy scale</td>
<td>12</td>
</tr>
<tr>
<td>2.3</td>
<td>Feynman diagram for the contribution at Leading Order to the Drell-Yan process</td>
<td>15</td>
</tr>
<tr>
<td>2.4</td>
<td>Feynman diagrams for examples of higher-order contributions to the Drell-Yan process</td>
<td>15</td>
</tr>
<tr>
<td>2.5</td>
<td>Schematic illustration of parton interactions within a proton</td>
<td>17</td>
</tr>
<tr>
<td>2.6</td>
<td>Proton NLO PDFs from the MSTW collaboration at two momentum transfer scales</td>
<td>18</td>
</tr>
<tr>
<td>2.7</td>
<td>Standard Model cross section predictions for proton-(anti)proton collisions</td>
<td>20</td>
</tr>
<tr>
<td>2.8</td>
<td>Summary of several Standard Model production cross section measurements performed by the ATLAS collaboration</td>
<td>21</td>
</tr>
<tr>
<td>2.9</td>
<td>The Bullet Cluster</td>
<td>23</td>
</tr>
<tr>
<td>4.1</td>
<td>The CERN accelerator complex</td>
<td>34</td>
</tr>
<tr>
<td>4.2</td>
<td>Magnets in the LHC tunnel</td>
<td>36</td>
</tr>
<tr>
<td>4.3</td>
<td>LHC magnet components</td>
<td>36</td>
</tr>
<tr>
<td>4.4</td>
<td>Placement of the dipole and quadrupole magnets responsible for bringing the LHC beams into collision at the ATLAS interaction point. Also shown are the positions of the LUCID, Zero Degree Calorimeter and ALFA detectors</td>
<td>37</td>
</tr>
<tr>
<td>4.5</td>
<td>Cumulative luminosity versus time delivered to ATLAS, recorded by ATLAS, and certified to be good quality data during stable beams and for proton-proton collisions at 7 and 8 TeV centre-of-mass energy in 2011 and 2012</td>
<td>38</td>
</tr>
<tr>
<td>4.6</td>
<td>Luminosity-weighted distribution of the mean number of interactions per crossing for the dataset collected in 2011 and 2012</td>
<td>38</td>
</tr>
<tr>
<td>4.7</td>
<td>$Z \rightarrow \mu^+\mu^-$ candidate event with reconstructed vertices from 25 simultaneous proton-proton interactions in the ATLAS detector</td>
<td>39</td>
</tr>
<tr>
<td>4.8</td>
<td>The ATLAS detector</td>
<td>41</td>
</tr>
</tbody>
</table>
Figure 5.10 Example fit results for MS towers with good angular resolution in data at $\sqrt{s} = 7\text{TeV}$.

Figure 5.11 Example fit results for Muon Spectrometer towers with known poor angular resolution in data at $\sqrt{s} = 7\text{TeV}$.

Figure 5.12 Example fit results for selected 2-station muons from data at $\sqrt{s} = 7\text{TeV}$, in momentum bins.

Figure 5.13 Momentum resolution of 2-station muons in data at $\sqrt{s} = 7\text{TeV}$.

Figure 5.14 Momentum resolution of 2-station muons in simulation at $\sqrt{s} = 7\text{TeV}$.

Figure 5.15 Momentum resolution of 2-station muons in simulation at $\sqrt{s} = 7\text{TeV}$ from comparing the refit 2-station momentum to MC truth.

Figure 6.1 Signal acceptance times efficiency for a $Z'$$_{SSM}$ boson as a function of $M_{Z'}$.

Figure 6.2 Yield per pb$^{-1}$ for each run, for the primary and secondary dimuon selection at $\sqrt{s} = 7\text{TeV}$.

Figure 6.3 Yield per pb$^{-1}$ for each run, for the primary and secondary dimuon selection at $\sqrt{s} = 8\text{TeV}$.

Figure 6.4 Event with the highest dimuon invariant mass observed in data collected by the ATLAS experiment in Run-I of the LHC.

Figure 6.5 Sketch of a proton-proton collision at high energy.

Figure 6.6 Perturbative QCD correction factor for $Z/\gamma^* \rightarrow \mu^+ \mu^-$ production as function of $m_{\mu^+ \mu^-}$ at $\sqrt{s} = 7\text{TeV}$.

Figure 6.7 Electroweak and photon-induced correction factor for $Z/\gamma^* \rightarrow \mu^+ \mu^-$ production as function of $m_{\mu^+ \mu^-}$ at $\sqrt{s} = 7\text{TeV}$.

Figure 6.8 Generated $Z$ $p_T$ spectra near the $Z$ resonance region and at higher invariant masses.

Figure 6.9 Perturbative QCD, electroweak and photon-induced correction factors for $Z/\gamma^* \rightarrow \mu^+ \mu^-$ production as function of $m_{\mu^+ \mu^-}$ at $\sqrt{s} = 8\text{TeV}$.

Figure 6.10 Muon track-based isolation distribution, immediately before the isolation cut.

Figure 6.11 Dimuon invariant mass distribution from events with two muons passing the reversed isolation requirement along with the rest of the selection.

Figure 6.12 Dimuon invariant mass distribution after final selection for the first 1.21 fb$^{-1}$ of data collected in 2011.

Figure 6.13 Cosmic muon event with hits in all Barrel detectors.

Figure 6.14 Example $Z'_{SSM}$ signal templates generated at $\sqrt{s} = 8\text{TeV}$.

Figure 6.15 Example 2D combined template of $Z/\gamma^*$ background and $Z'_{\text{Min}}$ signal generated at $\sqrt{s} = 8\text{TeV}$, as a function of $\gamma^*$ and $m_{\mu^+ \mu^-}$, for $M_{Z'_{\text{Min}}} = 2.5\text{TeV}$ and $\theta_{\text{Min}} = 0$.

Figure 6.16 Dimuon invariant mass in the selected events at $\sqrt{s} = 7\text{TeV}$.

Figure 6.17 Dimuon invariant mass in the selected events at $\sqrt{s} = 8\text{TeV}$.
Figure 6.18 Dimuon transverse momentum and rapidity in the selected events at $\sqrt{s} = 7\text{TeV}$ 108
Figure 6.19 Dimuon transverse momentum and rapidity in the selected events at $\sqrt{s} = 8\text{TeV}$ 109
Figure 6.20 Dimuon transverse momentum in the first 1.21 fb$^{-1}$ of data collected at $\sqrt{s} = 7\text{TeV}$ 110
Figure 6.21 Transverse momentum of the leading muon and sub-leading muon in the selected events at $\sqrt{s} = 7\text{TeV}$ 111
Figure 6.22 Transverse momentum of the leading muon and sub-leading muon in the selected events at $\sqrt{s} = 8\text{TeV}$ 112
Figure 6.23 Muon transverse momentum in the first 1.21 fb$^{-1}$ of data collected at $\sqrt{s} = 7\text{TeV}$ 113
Figure 6.24 $\eta$ and $\phi$ distributions for the selected muons at $\sqrt{s} = 7\text{TeV}$ 114
Figure 6.25 $\eta$ distributions for the leading muon and sub-leading muon in the selected events at $\sqrt{s} = 8\text{TeV}$ 115
Figure 6.26 $\phi$ distributions for the leading muon and sub-leading muon in the selected events at $\sqrt{s} = 8\text{TeV}$ 116
Figure 6.27 Missing transverse energy in the selected events at $\sqrt{s} = 7\text{TeV}$. 117
Figure 6.28 Two-dimensional histograms of $E_T^{\text{miss}}$ vs. $m_{\mu^+\mu^-}$ and $E_T^{\text{miss}}$ vs. leading muon $p_T$ for events passing the primary and secondary dimuon selection in collision data at $\sqrt{s} = 7\text{TeV}$. 118
Figure 6.29 Missing transverse energy in the selected events at $\sqrt{s} = 8\text{TeV}$, for the primary and secondary dimuon selection. 119
Figure 6.30 Two-dimensional histograms of $E_T^{\text{miss}}$ vs. $m_{\mu^+\mu^-}$ and $E_T^{\text{miss}}$ vs. leading muon $p_T$ for events passing the primary and secondary dimuon selection in collision data at $\sqrt{s} = 8\text{TeV}$. 120

Figure 7.1 Fractional uncertainty at 90% CL on the quark-antiquark luminosity at $\sqrt{s} = 7\text{TeV}$ due to PDFs 123
Figure 7.2 Stopping power for positive muons in copper as a function of $\beta\gamma = p/Mc$ . 125
Figure 7.3 Muon critical energy for the chemical elements . 125
Figure 7.4 Uncertainty on the background estimate due to the muon resolution as a function of $m_{\mu^+\mu^-}$ at $\sqrt{s} = 7\text{TeV}$ . 126
Figure 7.5 $Z'$ signal templates at 2 and 3 TeV for the primary dimuon selection at $\sqrt{s} = 7\text{TeV}$, with nominal smearing and over-smearing increasing the MS resolution smearing constants by their uncertainty . 126
Figure 7.6 Asymmetric uncertainty on the $Z/\gamma^*$ cross section at $\sqrt{s} = 8\text{TeV}$ as a function of $m_{\mu^+\mu^-}$ due to each PDF eigenvector taken separately. Here eigenvectors 1 to 8 are shown. 130
Figure 7.7 Asymmetric uncertainty on the $Z/\gamma^*$ cross section at $\sqrt{s} = 8$TeV as a function of $m_{\mu^+\mu^-}$ due to each PDF eigenvector taken separately. Here eigenvectors 9 to 16 are shown. ................................................................. [131]

Figure 7.8 Asymmetric uncertainty on the $Z/\gamma^*$ cross section at $\sqrt{s} = 8$TeV as a function of $m_{\mu^+\mu^-}$ due to each PDF eigenvector taken separately. Here eigenvectors 17 to 20 are shown. ................................................................. [132]

Figure 7.9 Asymmetric uncertainty on the $Z/\gamma^*$ cross section at $\sqrt{s} = 8$TeV as a function of the dilepton invariant mass due to the four distinct PDF eigenvector groups. ................................................................. [134]

Figure 7.10 Symmetric uncertainty on the $Z/\gamma^*$ cross section at $\sqrt{s} = 8$TeV as a function of the dilepton invariant mass resulting from the addition in quadrature of the uncertainties from the four PDF eigenvector groups. ................................................................. [135]

Figure 7.11 Symmetric uncertainty on the $Z/\gamma^*$ cross section at $\sqrt{s} = 8$TeV as a function of the dilepton invariant mass obtained using the MSTW prescription calculated with VRAP, shown along with the uncertainties due to the QCD scale, $\alpha_S$ variations, photon-induced corrections and higher-order electroweak corrections, as well as the difference between the cross section central values from ABM11 and the PDF variation uncertainty envelope from MSTW, taken as an additional systematic uncertainty. ................................................................. [136]

Figure 8.1 Differences between data and expectation in the and dimuon and dielectron channels at $\sqrt{s} = 7$TeV, with both the statistical and systematic uncertainties taken into account to derive a bin-by-bin local significance. ................................................................. [140]

Figure 8.2 Differences between data and expectation in the and dimuon and dielectron channels at $\sqrt{s} = 8$TeV, with both the statistical and systematic uncertainties taken into account to derive a bin-by-bin local significance. ................................................................. [141]

Figure 8.3 Absolute value of the LLR used in the search, as a function of the $Z'$ signal mass $M_{Z'}$ and cross section $\sigma_{Z'}$, using the $\sqrt{s} = 7$TeV dataset. ................................................................. [144]

Figure 8.4 Distribution of the most signal-like LLR found in each of 10,000 pseudo-experiments at $\sqrt{s} = 8$TeV. ................................................................. [145]

Figure 8.5 Observed and expected upper limits at 95% CL on $\sigma B$ for $Z'_{\text{SSM}}$ and $E_6 Z'$ boson production at $\sqrt{s} = 7$TeV. ................................................................. [147]

Figure 8.6 Observed and expected upper limits at 95% CL on $\sigma B$ for $E_6 Z'_\psi$ production at $\sqrt{s} = 8$TeV for the combination of the dielectron and dimuon channels. ................................................................. [148]

Figure 8.7 Observed and median expected upper limits at 95% CL on $\sigma B$ for $Z'_{\text{SSM}}$ production at $\sqrt{s} = 8$TeV for the exclusive dimuon and dielectron channels, and for both channels combined. ................................................................. [149]
Figure 8.8 Ratio of the observed limits at 95% CL for the $Z'_{SSM}$ search to the $Z'_{SSM}$ cross section times branching fraction for the combination of dielectron and dimuon channels ................................................................. 149
Figure 8.9 Observed and expected limits at 95% CL on $\gamma'$ as a function of the $Z'$ mass at $\sqrt{s} = 7\text{TeV}$ ................................................................. 151
Figure 8.10 Observed and expected limits at 95% CL on $\gamma'$ as a function of the $Z'$ mass at $\sqrt{s} = 8\text{TeV}$ ................................................................. 152
Figure 8.11 Observed and expected limits at 95% CL on $\gamma'$ as a function of $\theta_{\text{Min}}$ at $\sqrt{s} = 8\text{TeV}$ ........................................ 153
Figure 9.1 Ratios of the parton luminosity accessible at the LHC at $\sqrt{s} = 13\text{TeV}$ compared to that at $\sqrt{s} = 8\text{TeV}$ for gluon-gluon, gluon-quark and quark-antiquark processes ........................................ 156
Figure 9.2 Expected upper limits at 95% CL on $\sigma B$ for $Z'_{SSM}$ boson production for the projected HL-LHC dataset in the dimuon channel ........................................ 156
Figure A.1 Higgs boson production cross section by channel as a function of $M_H$ ......................................................... 173
Figure A.2 Higgs boson cross section times branching fraction to observable final states as a function of $M_H$ ................................................................. 173
Figure A.3 Histograms with data from the ATLAS experiment displaying evidence for the production of a Higgs boson in four different channels ........................................ 174
Figure A.4 Feynman diagram for VBF $H \rightarrow WW^* \rightarrow \ell\ell\nu\ell\nu$ at Leading Order ........................................ 175
Figure A.5 Transverse mass $m_T$ used as discriminant in the final stage of the cut-based search for VBF $H \rightarrow WW^* \rightarrow \ell\ell\nu\ell\nu$ ........................................ 176
Figure A.6 Light-quark jet rejection factor as a function of the $b$-tagging efficiency for different $b$-tagging algorithms ........................................ 178
Figure A.7 Possible spin configurations following a $H \rightarrow WW^* \rightarrow \ell\ell\nu\ell\nu$ decay where the $W$ bosons have non-zero spin in the direction of motion ........................................ 180
Figure A.8 VBF $H \rightarrow WW^* \rightarrow \ell\ell\nu\ell\nu$ candidate event ........................................ 181
Figure A.9 Schematic view of a simple decision tree ........................................ 182
Figure A.10 Illustration of the differences between a cut-based selection and one using a decision tree ........................................ 183
Figure A.11 Schematic diagram of the analysis fit model ........................................ 186
Figure A.12 Schematic diagram of the BDT optimization algorithm ........................................ 188
Figure A.13 Histograms of the BDT score in the Top CR for the DF channel and the SF channel ........................................ 189
Figure A.14 Histograms of the BDT training variables in the Top CR for the DF channel ........................................ 190
Figure A.15 Histograms of the BDT training variables in the Top CR for the SF channel ........................................ 191
Figure A.16 Histograms of the BDT training variables in the low-BDT VR for the DF channel ........................................ 192
Figure A.17 Histograms of the BDT training variables in the low-BDT VR for the SF channel ........................................ 193
Figure A.18 Distribution of $\rho(m_{\ell\ell},m_T)$ in the low-BDT VR ........................................ 194
Figure A.19 Example pair of 2D correlation histograms ........................................... 195
Figure A.20 Distributions of $m_{\ell\ell}$ and $m_{\tau\tau}$ in the $Z/\gamma' \to \tau^+\tau^-$ CR for the combination of the DF and SF channels .......................................................... 196
Figure A.21 Distributions of BDT score in the SF channel, for the low-$E_T^{\text{miss}}$ Z CR, region C and region D .......................................................... 198
Figure A.22 Illustration of the absence of correlation between the BDT score and $E_T^{\text{miss}}$ in simulated events .......................................................... 199
Figure A.23 Example object systematic uncertainty, due to the $p_T$, track resolution, on the BDT shape in the SR of the DF channel, for signal and the sum of all backgrounds 200
Figure A.24 Illustration of the sampling algorithm to generate pseudo-experiments ........ 201
Figure A.25 Number of events in pseudo-datasets after the pre-selection of the BDT analysis and the selection of the cut-based analysis ................................................. 202
Figure A.26 Signal significance and strength observed in 200 pseudo-experiments, with a BDT score as the discriminating variable ................................................. 203
Figure A.27 Signal significance and strength observed in 200 pseudo-experiments, with $m_T$ as the discriminating variable ................................................. 203
Figure A.28 Differences in signal significance and strength observed in individual pseudo-experiments between the two analyses ................................................. 204
Figure A.29 Two-dimensional histograms of the signal significance and strength observed in individual pseudo-experiments in the two analyses ................................................. 205
Figure A.30 Linearity test between the two analyses, for the signal significance and signal strength .......................................................... 205
Figure A.31 Event-by-event probabilities calculated in a previous definition of the DF Top CR under the hypotheses of $t\bar{t}$, WW, gluon-fusion Higgs, $Z/\gamma'$ and VBF Higgs production .......................................................... 207
Figure A.32 Event-by-event probabilities calculated in a previous definition of the blinded DF SR under the hypotheses of $t\bar{t}$, WW, gluon-fusion Higgs, $Z/\gamma'$ and VBF Higgs production .......................................................... 208
Figure A.33 Measurements of the signal strength ratios between the bosonic and fermionic Higgs production channels for the individual final states and their combination. 210
Figure A.34 Fit results for a parametrization of Higgs boson coupling strengths probing different scale factors for fermions and bosons. .......................................................... 210
<table>
<thead>
<tr>
<th>Acronym</th>
<th>Description</th>
</tr>
</thead>
<tbody>
<tr>
<td>BAT</td>
<td>Bayesian Analysis Toolkit</td>
</tr>
<tr>
<td>BDT</td>
<td>Boosted Decision Tree</td>
</tr>
<tr>
<td>BEE</td>
<td>Barrel Endcap Extra</td>
</tr>
<tr>
<td>BI</td>
<td>Barrel Inner</td>
</tr>
<tr>
<td>BM</td>
<td>Barrel Middle</td>
</tr>
<tr>
<td>BO</td>
<td>Barrel Outer</td>
</tr>
<tr>
<td>CKM</td>
<td>Cabbibo-Kobayashi-Maskawa</td>
</tr>
<tr>
<td>CL</td>
<td>Confidence Level</td>
</tr>
<tr>
<td>CP</td>
<td>Charge-Parity</td>
</tr>
<tr>
<td>CR</td>
<td>Control Region</td>
</tr>
<tr>
<td>CSC</td>
<td>Cathode Strip Chamber</td>
</tr>
<tr>
<td>DAQ</td>
<td>Data Acquisition</td>
</tr>
<tr>
<td>DCS</td>
<td>Detector Control System</td>
</tr>
<tr>
<td>DF</td>
<td>Different-Flavour</td>
</tr>
<tr>
<td>EI</td>
<td>Endcap Inner</td>
</tr>
<tr>
<td>EE</td>
<td>Endcap Extra</td>
</tr>
<tr>
<td>EM</td>
<td>Endcap Middle</td>
</tr>
<tr>
<td>EO</td>
<td>Endcap Outer</td>
</tr>
<tr>
<td>EPD</td>
<td>Event Probability Discriminant</td>
</tr>
<tr>
<td>Abbreviation</td>
<td>Description</td>
</tr>
<tr>
<td>--------------</td>
<td>------------------------------------</td>
</tr>
<tr>
<td>GUT</td>
<td>Grand Unified Theories</td>
</tr>
<tr>
<td>HL-LHC</td>
<td>High-Luminosity LHC</td>
</tr>
<tr>
<td>ID</td>
<td>Inner Detector</td>
</tr>
<tr>
<td>LEP</td>
<td>Large Electron-Positron Collider</td>
</tr>
<tr>
<td>LHC</td>
<td>Large Hadron Collider</td>
</tr>
<tr>
<td>LLR</td>
<td>Log-Likelihood Ratio</td>
</tr>
<tr>
<td>LO</td>
<td>Leading Order</td>
</tr>
<tr>
<td>MC</td>
<td>Monte Carlo</td>
</tr>
<tr>
<td>MDT</td>
<td>Muon Drift Tube</td>
</tr>
<tr>
<td>ME</td>
<td>Matrix Element</td>
</tr>
<tr>
<td>MS</td>
<td>Muon Spectrometer</td>
</tr>
<tr>
<td>NF</td>
<td>Normalization Factor</td>
</tr>
<tr>
<td>NLO</td>
<td>Next-to-Leading Order</td>
</tr>
<tr>
<td>NNLO</td>
<td>Next-to-Next-to-Leading Order</td>
</tr>
<tr>
<td>QCD</td>
<td>Quantum Chromodynamics</td>
</tr>
<tr>
<td>QED</td>
<td>Quantum Electrodynamics</td>
</tr>
<tr>
<td>PMNS</td>
<td>Pontecorvo-Maki-Nakagawa-Sakata</td>
</tr>
<tr>
<td>PDF</td>
<td>Parton Distribution Function</td>
</tr>
<tr>
<td>RPC</td>
<td>Resistive Plate Chamber</td>
</tr>
<tr>
<td>SCT</td>
<td>Semiconductor Tracker</td>
</tr>
<tr>
<td>SF</td>
<td>Same-Flavour</td>
</tr>
<tr>
<td>SM</td>
<td>Standard Model</td>
</tr>
<tr>
<td>SR</td>
<td>Signal Region</td>
</tr>
<tr>
<td>SSM</td>
<td>Sequential Standard Model</td>
</tr>
<tr>
<td>TGC</td>
<td>Thin Gap Chamber</td>
</tr>
<tr>
<td>Acronym</td>
<td>Description</td>
</tr>
<tr>
<td>---------</td>
<td>------------------------------</td>
</tr>
<tr>
<td>TMVA</td>
<td>Toolkit for Multivariate Data Analysis</td>
</tr>
<tr>
<td>TRT</td>
<td>Transition Radiation Tracker</td>
</tr>
<tr>
<td>VBF</td>
<td>Vector Boson Fusion</td>
</tr>
<tr>
<td>VR</td>
<td>Validation Region</td>
</tr>
</tbody>
</table>
Acknowledgements

My journey as a graduate student at UBC has been a wonderful one, not least due to the people I have met along the way, both in beautiful Vancouver and abroad.

First of all, I want to thank my supervisors, Colin Gay and Oliver Stelzer-Chilton, for having been outstanding mentors. You have provided me with sound advice and strong support when I needed it, as well as flexibility and freedom regarding research projects and travel plans, and have taught me the scientific knowledge and skills necessary to carry out cutting-edge research in particle physics. Particularly to Oliver who has been my closest collaborator in the last 5 years: thank you so much for having made time to meet and review results and documents, sometimes on short notice when the situation required it. Few students have had the chance to have such a positive relationship with their supervisors.

Thanks to Ashutosh Kotwal for your pragmatic advice, which taught me to concentrate my attention on non-negligible effects, and for sharing your insights in both theoretical and experimental domains. Thanks to Fabienne Ledroit for your attention to detail during the analysis and documentation stages: this helped tremendously, both to speed up the approval process at the time, and also to write this dissertation years later. Thanks to Bernd Stelzer for your help with statistics and analysis methods, and for forming such a great duo with your twin brother.

Thanks to Stephane Willocq, Dominique Fortin and all the other members of the very-high-momentum muon task force for your support as we worked towards understanding this crucial component of our detector’s performance.

To all collaborators from the Z’ analysis group, thank you for having been such a motivated, dedicated and strong team. It has been a pleasure to work with you, and may our paths cross again! Thanks also to everyone who contributed to writing or reviewing our internal and public documentation of the analysis. Special thanks to Kathy Copic and Wojtek Fedorko for your extraordinary leadership towards our first publications, and to Sarah Heim for having been such a fantastic teammate in analysis and convenership. Thanks also to Kelsey Allen for having been an amazing research assistant on this project.

Thanks to collaborators who provided me with example computer programs to get me started on projects, in particular to Ian Nugent and Doug Schouten for analysis frameworks. I am very grateful
for not having had to “re-invent the wheel”, allowing for faster progress in my research, and for the opportunity to learn by osmosis from your programming skills.

Thanks to my colleagues participating on the Higgs to WW∗ analysis, especially VBF analysts and multivariate analysis developers in particular: as the motto goes, “work hard, party hard”!

More generally, thanks to the ATLAS collaboration and the LHC at large: this is true teamwork. It is an honour to be part of such an awesome international endeavour.

Thanks to all my friends at UBC, TRIUMF and elsewhere for bringing sunshine in the workplace. I also want to thank UBC and TRIUMF for providing me the resources needed to carry out my graduate program, and for being such lively working environments. Thanks to my supervisory committee members: Mark Halpern, Christopher Hearty and Gordon Semenoff.

I acknowledge financial support from the Vanier Canada Graduate Scholarship program and the Natural Sciences and Engineering Research Council of Canada. I am grateful for generous scholarships and also for the ATLAS Canada travel budget, which allowed me to travel to CERN multiple times every year and to participate in several academic conferences worldwide. These travels have enabled me to actively participate in the vibrant particle physics community.

During my first three years in Vancouver, I have had the privilege of living at Green College, an outstanding community at UBC where I formed close friendships, and even met my wife! Thanks to Mark Vessey for your words of wisdom, wit and support, and may Green College keep blossoming.

To my friends here: I sincerely hope that we will keep in touch after scattering from Vancouver in all directions. Here is to more board games, potlucks, pub nights and hikes!


À mes parents, Johanne et Janot : merci, merci pour tout. Vous avez toujours eu une confiance absolue en mes aptitudes et m’avez soutenu à travers toutes les épreuves que j’ai rencontrées. Vous êtes les meilleurs parents du monde et je sais que vous serez également des grands-parents extraordinaires !

Francis, Julie et Jean-Michel, je suis super chanceux de vous avoir comme frères et soeur, et j’aimerais être en mesure de passer plus de temps avec vous.

Yuan, ma chérie !!! Thank you for your patience and support during my research and the preparation of my thesis. Especially during the many months I had to spend away from you traveling, I realized how truly special our love is. As much as I enjoy exploring the world, it is never the same without you. You make each hour of our everyday life at home so much more meaningful. Your love and care mean the world to me. I love you more than ever, and am so excited to start our new life as parents in the next months! And let’s hope that you will have as much fun reading my dissertation as I will with yours 😊
À ma famille :
passée, présente et future
Chapter 1

Introduction

The Standard Model of particle physics [108, 150, 160, 168] is the most successful theory of
elementary particle physics so far. Its accuracy in describing physical phenomena in an extremely
wide range of contexts is unparalleled, from the structure of atoms and molecules to the inner work-
ings of their constituents. Experiments in high-energy particle physics represent another domain
where the predictions of the Standard Model are verified. Building on the success of the CDF and
D0 experiments at the Tevatron, where the top quark was discovered in 1995 [65, 89], the ATLAS
and CMS experiments at the Large Hadron Collider (LHC) are now at the forefront of the energy
frontier. The discovery in 2012 of a Higgs boson [23, 76], the last particle predicted by the Standard
Model, is a momentous achievement.

Nevertheless, the Standard Model’s description of nature is incomplete. First of all, it neglects
gravitational effects: a more fundamental theory is necessary to reconcile general relativity with
quantum field theory. Additional theoretical problems such as the hierarchy of energy scales and
the large number of postulated parameters need to be addressed, while on the experimental side,
the observation of dark matter, dark energy and neutrino masses are clear indications that there is
more to be discovered. For this reason, high-energy particle physics experiments are designed to be
sensitive to a wide range of new physics, exploring uncharted territory in hopes of detecting new
phenomena that reach beyond the Standard Model.

In particular, searches for resonances decaying into electron and muon pairs have a long and
fruitful history [115]. Indeed, the discoveries in these channels of the J/ψ meson [37, 38], the
Y meson [116] and the Z boson [162, 164] all represent major breakthroughs in particle physics.
The main advantage of this search signature are that the signal is fully reconstructible and forms a
peak above relatively low and well-understood backgrounds. These features and the fact that many
hypotheses beyond the Standard Model predict the existence of such new resonances above the Z
mass peak make the search for new physics in this channel especially promising.

This dissertation describes searches for new neutral high-mass resonances decaying into muon
pairs using the full proton-proton collision dataset collected by the ATLAS experiment in Run-I of the LHC, including 5.0 fb$^{-1}$ collected at a centre-of-mass energy of 7 TeV [25] and 20.5 fb$^{-1}$ collected at a centre-of-mass energy of 8 TeV [36]. Chapter 2 describes the Standard Model of particle physics in the framework of quantum field theory, and discusses its successes and limitations. Chapter 3 presents hypotheses that reach beyond the Standard Model, with particular emphasis on models postulating new gauge symmetries. The existence of these additional symmetries often implies the presence of new particles to which this analysis could be sensitive. In Chapter 4 the LHC and the ATLAS detector are described. The performance of the ATLAS detector for reconstructing muons at very high momentum is assessed in Chapter 5, which also motivates the dedicated selection criteria that are employed. The next chapters describe the analysis itself: Chapter 6 describes the event selection as well as the techniques used to predict background and signal yields, and presents a comparison of the background expectations to data, Chapter 7 discusses systematic uncertainties and Chapter 8 presents the statistical methods used and the results of the search. Finally, Chapter 9 explores future directions and Chapter 10 concludes with a summary.

In parallel with the work presented in this dissertation, significant contributions were also made to the multivariate analysis looking for Vector Boson Fusion production of Higgs bosons decaying into W boson pairs in final states with two electrons or muons. Measuring the production rate of the newly-discovered Higgs boson in this channel will provide one of the best constraints on the Higgs coupling to W bosons. Efforts toward completing the results of this analysis are still ongoing at the time of writing this dissertation, preventing their inclusion in this dissertation. Nevertheless, a summary of contributions to this Higgs boson analysis is presented in Appendix A.
Chapter 2

The Standard Model of Particle Physics

This chapter presents the Standard Model of particle physics, which is so far the most successful theory describing non-gravitational interactions. Instead of a historical presentation where experimental discoveries and theoretical advances are presented chronologically, a more axiomatic format is followed\(^1\). Such a format allows for a more logical presentation of the theory, although it does not present the crucial experiments and early concepts that informed its development\(^2\).

First, the Lagrangian of the Standard Model is presented in Section 2.1 and a few examples of its successful predictions are discussed in Section 2.2. Section 2.3 describes the theoretical and experimental shortcomings of the Standard Model, which are indicative of the challenges that hypotheses going beyond it have to meet.

2.1 The Standard Model Lagrangian

The Standard Model of particle physics is a renormalizable quantum field theory with local gauge invariance. The gauge symmetry is postulated to be \(SU(3)_C \times SU(2)_L \times U(1)_Y\), where the subscripts represent the conserved quantum numbers associated with each symmetry group. The conserved quantum numbers corresponding to \(SU(3)_C\) are called “colours”, and the gauge bosons are named “gluons”. \(SU(2)_L\) symmetry only interacts with left-handed fermions and their right-handed antiparticles, and the conserved quantity is the “isospin” \(T_3\). Finally, the conserved quantum number of the \(U(1)_Y\) symmetry is the “hypercharge” \(Y\).

2.1.1 Forces and Matter

As shown in Table 2.1, the vector bosons of the Standard Model correspond to the adjoint representations of \(SU(3)_C\) and \(SU(2)_L\), with an additional vector field corresponding to \(U(1)_Y\) which

---

\(^1\) This chapter borrows many elements from Ref. [58] and from lecture notes by J. Ng. (unpublished).

The metric convention used is \(\eta_{\mu\nu} = \text{diag}(-, +, +, +)\), the operator \(\partial_\mu \equiv \partial / \partial x^\mu\), and \(\gamma^\mu\) are Dirac matrices.

\(^2\) A captivating historical narrative can be found in Ref. [73].
transforms as a singlet. The notation \((C, L, Y)\) is used to describe how fields transform under the Standard Model gauge symmetry group: specifically, \(C\) gives the corresponding representation of the \(SU(3)_C\) algebra, and similarly \(L\) for the \(SU(2)_L\) algebra.

**Table 2.1:** Gauge bosons in the Standard Model.

<table>
<thead>
<tr>
<th>Vector field</th>
<th>(SU(3)_C)</th>
<th>(SU(2)_L)</th>
<th>(U(1)_Y)</th>
</tr>
</thead>
<tbody>
<tr>
<td>(G^\alpha_\mu), (\alpha \in [1..8])</td>
<td>8</td>
<td>1</td>
<td>0</td>
</tr>
<tr>
<td>(W^a_\mu), (a \in [1..3])</td>
<td>1</td>
<td>3</td>
<td>0</td>
</tr>
<tr>
<td>(B_\mu)</td>
<td>1</td>
<td>1</td>
<td>0</td>
</tr>
</tbody>
</table>

**Table 2.2:** Fermions in the Standard Model, in flavour basis. Colour indices \(c \in \{r, g, b\}\) for quarks and generation indices \(m \in [1..3]\) for all fermions are omitted.

<table>
<thead>
<tr>
<th>Fermion field</th>
<th>(SU(3)_C)</th>
<th>(SU(2)_L)</th>
<th>(U(1)_Y)</th>
</tr>
</thead>
<tbody>
<tr>
<td>(U_L)</td>
<td>3</td>
<td>2</td>
<td>(1/6)</td>
</tr>
<tr>
<td>(D_L)</td>
<td>(3)</td>
<td>1</td>
<td>(2/3)</td>
</tr>
<tr>
<td>(U_R)</td>
<td>3</td>
<td>1</td>
<td>(-1/3)</td>
</tr>
<tr>
<td>(D_R)</td>
<td>(3)</td>
<td>1</td>
<td>(-1/3)</td>
</tr>
<tr>
<td>(v_L)</td>
<td>1</td>
<td>2</td>
<td>(-1/2)</td>
</tr>
<tr>
<td>(\ell_L)</td>
<td>(1)</td>
<td>1</td>
<td>(-1)</td>
</tr>
</tbody>
</table>

The Standard Model fermions come in three generations, and are collectively represented by \(\psi_m\), \(m \in [1..3]\). Fermions in the second and third generations interact with gauge bosons in exactly the same way as their counterparts in the first generation. They transform as fundamental representations of the gauge groups or as singlets, as detailed in Table 2.2. The fermions that interact
with $SU(3)_C$ are called “quarks”, and the ones that do not are called “leptons”. Fermions with left-handed helicity, $\psi_{m,L} = \frac{1}{2}(1 + \gamma^5)\psi_m$ where $\gamma^5 = -i\gamma^0\gamma^1\gamma^2\gamma^3$, transform differently from fermions with right-handed helicity, $\psi_{m,R} = \frac{1}{2}(1 - \gamma^5)\psi_m$. Specifically, the left-handed fermions are grouped together into $SU(2)_L$ doublets, with the ones on top having isospin $T_3 = 1/2$, and the ones on the bottom having isospin $T_3 = -1/2$. Right-handed fermions are assigned $T_3 = 0$.

The Lagrangian consisting of all renormalizable and gauge-invariant terms with the Standard Model fermion and gauge boson fields is the following:

$$\mathcal{L}_G = -\frac{1}{4} G^\alpha_{\mu\nu} G^{\alpha\mu\nu} - \frac{1}{4} W^a_{\mu\nu} W^{a\mu\nu} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu} - \frac{i}{2} (\bar{\psi}_m \gamma^\mu D_\mu \psi_m + \text{h.c.})$$

(2.1)

where the covariant field strengths are

$$G^\alpha_{\mu\nu} = \partial_\mu G^\alpha_\nu - \partial_\nu G^\alpha_\mu + g_3 f^\alpha_{\beta\gamma} G^\beta_\mu G^\gamma_\nu$$

(2.2)

$$W^a_{\mu\nu} = \partial_\mu W^a_\nu - \partial_\nu W^a_\mu + g_2 \varepsilon_{abc} W^b_\mu W^c_\nu$$

(2.3)

$$B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu$$

(2.4)

and the covariant derivative of the fermion fields is

$$D_\mu \psi_m = (\partial_\mu - ig_1 Y B_\mu - \frac{ig_2}{2} \sigma_\mu W^a_\mu - \frac{ig_3}{2} \lambda_\alpha G^\alpha_\mu) \psi_m$$

(2.5)

where the $W^a_\mu$ term only applies to left-handed fermions and right-handed antifermions, and the $G^\alpha_\mu$ term only applies to quarks. Here the $\sigma_\mu$ are Pauli matrices, the generators of $SU(2)$, and the $\lambda_\alpha$ are Gell-Mann matrices, the generators of $SU(3)$. The structure constants of $SU(3)$ are represented by $f^\alpha_{\beta\gamma}$, and those of $SU(2)$ are simply the fully antisymmetric $\varepsilon_{abc}$. The couplings $g_1$, $g_2$, $g_3 \in \mathbb{R}$ determine the strength of the interaction terms, and are discussed in more detail in Section 2.2.1.

This Lagrangian has a rather surprising feature: because of the absence of left-right symmetry in the Standard Model, no mass term is allowed for any of the fields. For instance, while mass terms for the fermions might have been written as

$$\mathcal{L}_M = -\frac{M_{mn}}{2} \bar{\psi}_m \psi_n$$

(2.6)

$$= -\frac{M_{mn}}{2} (\bar{\psi}_{m,L} \psi_{n,L} + \bar{\psi}_{m,R} \psi_{n,R})$$

these terms are not gauge-invariant under $U(1)_Y$: in other words there is no combination of fields $\bar{\psi}_{m,L} \psi_{n,R}$ or $\bar{\psi}_{m,R} \psi_{n,L}$ that is hypercharge-neutral. Direct mass terms for vector bosons are also forbidden by gauge invariance. Since we observe that most fundamental particles do have an intrinsic mass, the absence of mass terms constitutes a significant theoretical problem.

3 with a caveat discussed in Section 2.3.3
2.1.2 The Englert-Brout-Higgs-Guralnik-Hagen-Kibble Mechanism

The solution to this problem is to add to the model a scalar field that retains a non-zero value in the ground state and thereby breaks electroweak symmetry \([96, 112, 117]\). The simplest scalar field that can be used for this purpose transforms as \((1, 2, 1/2)\):

\[
\varphi = \begin{bmatrix} \varphi^+ \\ \varphi^0 \end{bmatrix}
\]  

where \(\varphi^+\) and \(\varphi^0\) are complex scalars. The potential of this scalar field can be parametrized as

\[
V(\varphi^\dagger \varphi) = \lambda (\varphi^\dagger \varphi - \mu^2 / 2\lambda)^2 \\
= \lambda (\varphi^\dagger \varphi)^2 - \mu^2 \varphi^\dagger \varphi + \mu^4 / 4\lambda
\]

with \(\lambda\) and \(\mu^2\) real and positive. This potential is minimized for \(\varphi^\dagger \varphi = \mu^2 / 2\lambda \equiv v^2 / 2\), where \(v\) is defined as the non-zero “vacuum expectation value” of the scalar field. A very useful form of \(\varphi\) is obtained by performing a gauge transformation such that no quadratic vector-scalar cross terms remain: this gauge choice is called “unitary gauge” and the scalar field becomes

\[
\varphi = \begin{bmatrix} 0 \\ \frac{1}{\sqrt{2}} (v + H) \end{bmatrix}
\]

with \(H\) a real scalar field named the Higgs boson. This gauge choice fixes the covariant derivative of the scalar field as follows:

\[
D_\mu \varphi = (\partial_\mu - \frac{ig_1}{2} B_\mu - \frac{ig_2}{2} \sigma_\mu W_\mu^a) \varphi
\]

\[
= \begin{bmatrix} -\frac{ig_2}{2\sqrt{2}} (W_\mu^1 - iW_\mu^2)(v + H) \\ \frac{1}{\sqrt{2}} \partial_\mu H - \frac{i}{2\sqrt{2}} (g_1 B_\mu - g_2 W_\mu^3)(v + H) \end{bmatrix}
\]

Then the Standard Model Lagrangian becomes:

\[
\mathcal{L}_{\text{SM}} = \mathcal{L}_{G\psi} + \mathcal{L}_H
\]
with
\[ \mathcal{L}_H = - (D_\mu \Phi)^\dagger (D^\mu \Phi) - V(\Phi^\dagger \Phi) - (y_{\ell, mn}[\bar{\nu}_{m,L} \bar{\ell}_{n,L}] \ell_{n,R} \phi) + y_{\varphi, mn}[\bar{\varphi}_{m,L} \varphi_{n,L}] \varphi_{n,R} \phi + y_{\varphi, mn}[\bar{\varphi}_{m,L} \varphi_{n,L}] \varphi_{n,R} \phi \]
\[ + \frac{1}{2} \partial_\mu H \partial^\mu H - \frac{\lambda}{4} v^H - \frac{\lambda}{2} |H|^4 \]
\[ - \frac{g_2^2}{8} (v + H)^2 |W_\mu^1 - i W_\mu^2|^2 - \frac{1}{8} (v + H)^2 (g_1 B_\mu - g_2 W_\mu^3)^2 \]
\[ - \frac{1}{\sqrt{2}} (v + H) (y_{\ell, mn} \bar{\ell}_{m,L} \ell_{n,R} + y_{\varphi, mn} \varphi_{m,L} \varphi_{n,R} + y_{\varphi, mn} \bar{\varphi}_{m,L} \varphi_{n,R} \phi + h.c.) \]

(2.13)

The complete Standard Model Lagrangian contains mass terms for most elementary particles. To see this more clearly, it is very useful to redefine the vector boson and fermion fields in the “mass basis”, in order to eliminate all quadratic mixed terms between fields. The vector boson fields are redefined as follows:
\[ W_\mu^\pm = \frac{1}{\sqrt{2}} (W_\mu^1 \mp i W_\mu^2) \]
\[ Z_\mu = \frac{g_2 W_\mu^3 - g_1 B_\mu}{\sqrt{g_1^2 + g_2^2}} \]
\[ A_\mu = \frac{g_1 W_\mu^3 + g_2 B_\mu}{\sqrt{g_1^2 + g_2^2}} \]
\[ \equiv W_\mu^3 \cos \theta_W - B_\mu \sin \theta_W \]
\[ \equiv W_\mu^3 \sin \theta_W + B_\mu \cos \theta_W \]

(2.14)

defining $\theta_W$, the weak mixing angle (or Weinberg angle). The fermion Yukawa terms are diagonalized using unitary matrices $U_{\psi}$, which mix the fermion fields from different generations:
\[ \bar{\psi}_{m,L} = U_{\psi, mn} \psi_{n,L} \]
\[ \bar{\psi}_{m,R} = U_{\psi, mn}^* \psi_{n,R} \]

(2.18)

such that $U_{\psi, mn}^* U_{\psi, mn} \equiv \delta_{mn}$ is diagonal.

These transformations simplify $\mathcal{L}_H$ considerably; in the mass basis, $\mathcal{L}_H$ becomes
\[ \mathcal{L}_H = - \partial_\mu H \partial^\mu H - \frac{\lambda}{4} v^H - \frac{\lambda}{2} |H|^4 \]
\[ - \frac{g_2^2}{4} (v + H)^2 |W_\mu^1 - W_\mu^2|^2 - \frac{1}{8} (v + H)^2 (g_1^2 + g_2^2) Z_\mu Z_\mu \]
\[ - \frac{y_{\psi, mn}}{\sqrt{2}} (v + H) \bar{\psi}_m \psi_m \]

(2.19)
yielding the following masses, all directly proportional to $v$:

$$M_W = \frac{g_2}{2} \cdot v \quad (2.20)$$

$$M_Z = \sqrt{\frac{g_1^2 + g_2^2}{2}} \cdot v \quad (2.21)$$

$$M_{\tilde{\psi}_m} = \frac{y_{\tilde{\psi}_m}}{\sqrt{2}} \cdot v \quad (2.22)$$

$$M_H = \sqrt{2\lambda} \cdot v \quad (2.23)$$

From the experimental values of $M_W$ and $M_Z$, we find $v = 246$ GeV.

Nine gauge bosons remain massless: the eight gluons $G^\alpha_{\mu}$, and $A_\mu$ which is identified as the photon. This implies that the scalar field $\varphi$ broke the symmetry of the Standard Model down from $SU(3)_C \times SU(2)_L \times U(1)_Y$ to the low-energy $SU(3)_C \times U(1)_Q$, where $Q = Y + T_3$ is identified as the electric charge. Gluons are often collectively represented by the letter $g$, and photons by the letter $\gamma$.

The standalone study of $U(1)_Q$ interactions is called Quantum Electrodynamics (QED), and that of $SU(3)_C$ interactions is called Quantum Chromodynamics (QCD) by analogy.

Neutrinos are also strictly massless in the Standard Model, this time in conflict with experimental results; this limitation is discussed in Section 2.3.4.

2.1.3 Electroweak Interactions Revisited

While the field redefinitions of Equations 2.15-2.18 simplify $\mathcal{L}_H$, this comes at the cost of a more intricate structure in $\mathcal{L}_{GW}$. Gluon self-interactions are unaffected, but $W$ and $Z$ bosons interact with photons in three electroweak self-interaction terms: $WW\gamma$, $WW\gamma\gamma$ and $WWZ\gamma$. Additional self-interaction terms $WWZ$, $WWZZ$ and $WWWW$ complete the list.

Further, while electroweak interactions with fermions were originally diagonal in fermion generation, as is manifest from Equation 2.11, the fermion field redefinitions of Equation 2.18 introduce terms in which a $W$ boson interacts with two fermions from different generations. In other words, while in the initial basis, called the “flavour basis”, all interactions between $W$ bosons and fermions preserve fermion generation but propagation in the vacuum does not (a phenomenon called fermion “flavour oscillation”), in the mass basis fermions are perceived as preserving their identity as they propagate in the vacuum but interactions with $W$ bosons may occur between fermions of different generations.

Because all quarks have sizeable masses, the mass basis is privileged in the quark sector. With increasing generation index $m$, the mass-basis quarks $u_m$ are named up ($u$), charm ($c$) and top ($t$) quarks, and the mass-basis quarks $d_m$ are named down ($d$), strange ($s$) and bottom ($b$) quarks. Quarks are collectively represented by the letter $q$. 
The interaction terms of quarks with $W$ bosons become

$$L_{Wq} = \frac{ig_2}{\sqrt{2}} (V_{mnL} \bar{q}_m \gamma^\mu p_n \gamma^\mu q_n + (V^\dagger)_{mnL} \bar{q}_m \gamma^\mu p_n \gamma^\mu q_n) \tag{2.24}$$

where $V = U^\dagger \nu U_d$ using the $U_\nu$ from Equation 2.18 is the Cabbibo-Kobayashi-Maskawa (CKM) matrix:

$$V = \begin{bmatrix}
V_{ud} & V_{us} & V_{ub} \\
V_{cd} & V_{cs} & V_{cb} \\
V_{td} & V_{ts} & V_{tb}
\end{bmatrix} = \begin{bmatrix}
1 - \frac{1}{2} \lambda^2 & \lambda & A \lambda^3 (\rho - i \eta) \\
-\lambda & 1 - \frac{1}{2} \lambda^2 & A \lambda^2 \\
A \lambda^3 (1 - \rho - i \eta) & -A \lambda^2 & 1
\end{bmatrix} + O(\lambda^4) \tag{2.25}
$$

The last line of Equation 2.25 defines the Wolfenstein parametrization of the CKM matrix, which is valid to fourth order in $\lambda \approx 0.225$, with $A \approx 0.8$, $\rho \approx 0.13$ and $\eta \approx 0.35$ [52]. The main advantage of this parametrization is that it highlights the relative size of the CKM matrix elements with respect to each other; in particular, the CKM matrix is nearly diagonal.

On the other hand, because neutrinos are massless in the Standard Model, it is conventional to conceptualize leptons in the flavour basis. The charged leptons $\ell_m$ are the electron $e$, the muon $\mu$ and the tauon $\tau$, each with an associated neutrino, labeled respectively $\nu_e$, $\nu_\mu$ and $\nu_\tau$. The discovery of neutrino flavour oscillations significantly enriched this description, as discussed in Section 2.3.4. The interaction terms of leptons with $W$ bosons are

$$L_{W\ell} = \frac{ig_2}{\sqrt{2}} (\bar{\nu}_m \gamma^\mu \ell_m \gamma^\mu q_n + \bar{\ell}_m \gamma^\mu \nu_m \gamma^\mu q_n) \tag{2.26}$$

Finally, the fermion interactions with the neutral electroweak gauge bosons are described by

$$L_{(Z/\gamma)\psi} = ig_{EM} \bar{\psi}_m \gamma^\mu Q \psi_m A_\mu + ig_Z \bar{\psi}_m \gamma^\mu \left( \frac{1 + \gamma^5}{2} T_3 - Q \sin^2 \theta_W \right) \psi_m Z_\mu \tag{2.27}$$

$$= ig_{EM} \bar{\psi}_m \gamma^\mu Q \psi_m A_\mu + ig_Z \bar{\psi}_m \gamma^\mu \left( \frac{1 + \gamma^5}{2} g_L + \frac{1 - \gamma^5}{2} g_R \right) \psi_m Z_\mu \tag{2.28}$$

where in the last expression, $g_L \equiv T_3 - Q \sin^2 \theta_W$ and $g_R \equiv -Q \sin^2 \theta_W$. The first term is identified with the Lagrangian of QED, with the factor $g_{EM} \equiv g_1 \cos \theta_W = g_2 \sin \theta_W$ identified as the electromagnetic coupling constant, i.e. the elementary charge. The coupling to the $Z$ boson is $g_Z \equiv g_1 / \sin \theta_W = g_2 / \cos \theta_W = \sqrt{g_1^2 + g_2^2} = 2M_Z / v$. It is interesting to notice that in contrast with interactions involving $W$ bosons, these neutral interactions cannot change the type or flavour of fermions.
2.1.4 Summary

The particle content of the Standard Model is encapsulated in Figure 2.1. In summary, it consists of eight gluons, mediators of the strong force, the photon, which mediates the electromagnetic force, and the massive weak bosons $W$ and $Z$, responsible for the weak force. Three generations of matter particles are present: quarks interact with each other via all three forces, charged leptons interact via both electroweak forces and neutrinos only interact via the weak force. Finally, the Higgs boson gives a mass to the quarks, charged leptons and weak bosons by interacting with all of them.

![Figure 2.1: Particle content of the Standard Model. Figure credit: Fermilab.](image)

2.2 Successes of the Standard Model

The experimental discovery of all the particles postulated by the Standard Model is an exceptional achievement, which certainly constitutes a major success of the theory. The latest discoveries of matter particles happened at Fermilab, where the top quark was observed in 1995 by the CDF and D0 experiments [65, 89], as well as the tau neutrino in 2000 by the DONUT experiment [95]. This endeavour culminated in the discovery of a Higgs boson in 2012 by the ATLAS and CMS experiments at the LHC [23, 76]. The experimental evidence so far indicates that this boson is compatible with expectations from the Standard Model\(^4\), thereby validating our understanding of the mechanism responsible for breaking the electroweak symmetry and for giving masses to elementary particles.

\(^4\) More details in Appendix A

10
But the Standard Model is much more than simply the sum of its parts. The next sections discuss two successful applications of the Standard Model, illustrating its ability to describe interactions between particles. While many more useful calculations could be illustrated, such as the lifetimes of unstable particles and the probabilities for each combination of their decay products, quantum corrections to the mass of particles, scattering amplitudes across different energy ranges, etc., the following two are chosen because of their relevance to the main topic of this dissertation.

First, implications of the specific values of the coupling strengths are discussed. These values are interpreted as originating from the evolution of the couplings with decreasing energy. They explain why quarks and gluons must combine into hadrons and why matter is organized in atoms at low energies. Then, the expected production rates for heavy particles at very-high energy colliders are discussed and compared with experimental results, with special attention given to the process $pp \rightarrow Z/\gamma^* \rightarrow \mu^+\mu^-$.

2.2.1 Asymptotic Freedom and Infrared Slavery

The larger size of the coupling $g_3$ with respect to $g_1$ and $g_2$ explains why the corresponding force is called the strong force. Indeed while at low energy scales the electroweak couplings $g_1, g_2 \ll 1$, which means that this part of the theory can be treated pertubatively, the coupling $g_3$ is significantly larger and, very importantly, gets larger towards low energies.

An equivalent way to express the couplings is in terms of

$$\alpha_i(\mu) \equiv \frac{g_i^2(\mu)}{4\pi}$$

where $\mu$ is the renormalization scale, which when performing calculations is taken to be close to the energy scale or momentum transfer scale of the process of interest. In particular, the strong coupling is $\alpha_S \equiv \alpha_3$:

$$\alpha_S(\mu) \equiv \frac{g_3^2(\mu)}{4\pi}$$

$$\alpha_S(M_Z) \approx 0.117$$

and the electromagnetic coupling is

$$\alpha_{EM}(\mu) \equiv \frac{g_{EM}^2(\mu)}{4\pi}$$

$$\alpha_{EM}(M_Z) \approx 1/128$$

The dependence of the gauge couplings of the Standard Model as a function of $\mu$, shown in
Figure 2.2, is described by the function

\[ \beta(\alpha) \equiv \mu^2 \frac{\partial \alpha}{\partial \mu^2} = \sum_{i=0}^{\infty} b_i \alpha^{i+2} \]  

(2.34)

Keeping only the first term of the sum, this is solved as

\[ \frac{1}{\alpha(\mu)} = \frac{1}{\alpha(\mu_0)} - b_0 \ln \frac{\mu^2}{\mu_0^2} \]  

(2.35)

where a convenient choice for \(\mu_0\) is \(M_Z\), as in the quoted values above. The coefficients \(b_0\) can be calculated explicitly, and the results for \(\alpha_S\) and \(\alpha_{\text{EM}}\), valid for \(\mu \gtrsim M_t\), are

\[ \beta(\alpha_S) \approx -\frac{7}{2\pi} \alpha_S^2 + \mathcal{O}(\alpha_S^3) < 0 \]  

(2.36)

\[ \beta(\alpha_{\text{EM}}) \approx \frac{2}{3\pi} \alpha_{\text{EM}}^3 + \mathcal{O}(\alpha_{\text{EM}}^4) > 0 \]  

(2.37)

Figure 2.2: Gauge couplings of the Standard Model as a function of energy scale [152].

It follows that in contrast with the strong coupling, the electromagnetic coupling grows with energy. This explains why while at very high energy scales \(\mu \sim 10^{14} \text{ GeV}\), called the “unification scale”, all couplings of the Standard Model have similar values, at low energies their strength varies importantly. Seen from another perspective, the strength of the electromagnetic force decreases as the distance between particles increases (as is familiar from the classical Coulomb law in \(1/r^2\)), but the effect of the strong force increases with the distance between quarks and gluons.
This implies that at very high energy, quarks and gluons essentially behave like free particles with respect to each other: this feature of the strong interaction is called “asymptotic freedom”. Its counterpart is “infrared slavery”: at low energy, gluons and quarks must group together into colourless bound states.

There are a few different possibilities for quarks to bind together. First, a quark and an antiquark can join to form a “meson”. The most common mesons are the pions, consisting of up and down quark-antiquark pairs. Other common mesons in high-energy particle physics are made of one up or down quark or antiquark, while the other antiquark or quark is of a heavier flavour: depending on this flavour such mesons are called kaons (strange), $D$-mesons (charm) and $B$-mesons (bottom). Charm quark-antiquark pairs called $J/\psi$ mesons, and bottom quark-antiquark pairs called $\Upsilon$ mesons are also particularly interesting because of their resonant production and decay into lepton pairs.

Alternatively, three quarks of different colours can join to form a “baryon”. Three antiquarks can also join similarly to form antibaryons. The most common baryons found in nature are the proton, made of two up quarks and one down quark, and the neutron, made of one up quark and two down quarks. The relative strengths of the couplings in the low-energy limit of the Standard Model therefore explains the structure of atoms, with quarks and gluons bound together in protons and neutrons forming atomic nuclei, and electrons in quantum orbits around these nuclei at effective distances inversely proportional to $\alpha_{EM}$.

2.2.2 Calculation of Cross Sections and Decay Rates

Perhaps the most important use of the Standard Model for physicists studying particle collisions at very-high energy is to calculate the rates for all the different possible outcomes of a collision: these rates are called the cross sections of each process of interest. This section first introduces concepts relevant when calculating cross sections, illustrated with an example, before comparing state-of-the-art calculations in the Standard Model to experimental measurements.

Scattering amplitudes and Feynman diagrams

In time-dependent perturbative quantum field theory, the amplitude connecting initial states $|\alpha\rangle$ to final states $|\beta\rangle$ is represented by the $S$-matrix:

$$S_{\beta\alpha} = \langle \beta | S | \alpha \rangle$$

from which it is conventional to factor out the interaction matrix element $\mathcal{M}$

$$S_{\beta\alpha} = \delta_{\beta\alpha} - i \mathcal{M}_{\beta\alpha} (2\pi)^4 \delta^4(k_{\beta} - p_{\alpha})$$
where \( p_\alpha \) represents the sum of the four-momenta of the initial-state particles, and \( k_\beta \) similarly for final-state particles. The link between the matrix element \( M \) and the Lagrangian density \( \mathcal{L} \) discussed earlier is via the interaction Hamiltonian density \( \mathcal{H}_I \):

\[
S = \sum_{n=0}^{\infty} \frac{(-i)^n}{n!} \int_{-\infty}^{\infty} d^4x_1 \ldots d^4x_n T[\mathcal{H}_I(x_1) \ldots \mathcal{H}_I(x_n)]
\]  

(2.40)

where \( T \) is the time-ordering operator, ensuring that the interactions happen in the proper sequence. The key point of the perturbative approach is that each factor of \( \mathcal{H}_I \) carries with it one factor of the coupling between particles; if these couplings are small compared to unity, then the first terms constitute a good approximation to the full amplitude. The matrix element \( M \) then becomes

\[
\mathcal{M}_{\beta\alpha} = \langle \beta | \mathcal{H}_I(0) | \alpha \rangle + \frac{(-i)^2}{2} \int d^4x \langle \beta | T[\mathcal{H}_I(x) \mathcal{H}_I(0)] | \alpha \rangle + \ldots
\]  

(2.41)

When the initial state consists of a single particle, this matrix element allows to calculate its differential decay rate \( d\Gamma \) as follows:

\[
d\Gamma = \frac{|\mathcal{M}_{\beta\alpha}|^2}{2E_\alpha} (2\pi)^4 \delta^4(k_\beta - p_\alpha) d\beta
\]  

(2.42)

\[
d\beta \equiv \prod_{f \in \beta} \frac{d^3k_f}{2E_f(2\pi)^3}
\]  

(2.43)

where \( E \) represents the energy of a particle, and \( k \) the three-momentum of final-state particles.

Similarly, in the context of two-body scattering, \( M \) appears in the expression for the differential cross section \( d\sigma \):

\[
d\sigma = \frac{|\mathcal{M}_{\beta\alpha}|^2}{4\sqrt{(p_1 \cdot p_2)^2 - m_1^2 m_2^2}} (2\pi)^4 \delta^4(k_\beta - p_\alpha) d\beta
\]  

(2.44)

where the subscripts 1 and 2 refer to the incoming particles in the initial state.

Feynman diagrams allow to conveniently calculate the matrix elements that enter such calculations in quantum field theory, in addition to providing a simple way to visualize interaction processes between particles. This technique is used in the following section to calculate the leading contribution to the Drell-Yan process at the LHC: \( pp \to Z/\gamma^* \to \mu^+ \mu^- \). The leading term in the series for the matrix element \( M \) in this case will involve two interactions: one between constituents of the incoming protons and a Z boson or a photon, and one between the same electroweak boson and the outgoing muons. Figure 2.3 shows the Feynman diagram corresponding to this term.

A calculation keeping only this first term in the series expansion for \( M \) is called a calculation at Leading Order (LO). Calculations keeping higher-order terms in the couplings are called calculations at Next-to-Leading Order (NLO), Next-to-Next-to-Leading Order (NNLO), and so on. Figure 2.4 shows examples of Feynman diagrams corresponding to higher-order contributions to the
Drell-Yan process. These additional contributions fall under two main categories: the emission of radiation contributing to the final state, and the presence of particle loops during the interaction.

Because each order of the perturbative series involves one more power of the interaction Hamiltonian and therefore of the coupling, corresponding to an additional vertex in the corresponding Feynman diagrams, the number of diagrams increases factorially, making the calculation of each successive term in \(M\) harder than the precedent. This becomes an issue in the context of processes involving the strong interaction, especially in the low-energy regime where the coupling \(\alpha_S\) is large. Fortunately, this is not a concern for the main processes of interest here.

![Feynman diagram for the contribution at Leading Order to the Drell-Yan process](image)

**Figure 2.3:** Feynman diagram for the contribution at Leading Order to the Drell-Yan process \(pp \rightarrow Z/\gamma^* \rightarrow \mu^+\mu^-\).

![Feynman diagrams for examples of higher-order contributions to the Drell-Yan process](image)

**Figure 2.4:** Feynman diagrams for examples of higher-order contributions to the Drell-Yan process \(pp \rightarrow Z/\gamma^* \rightarrow \mu^+\mu^-\): gluon emission in the initial state and photon emission in the final state (left), and a gluon loop in the initial state (right).

**The Drell-Yan process at the LHC:** \(pp \rightarrow Z/\gamma^* \rightarrow \mu^+\mu^-\)

As illustrated in Figure 2.3, the leading contribution to the process \(pp \rightarrow Z/\gamma^* \rightarrow \mu^+\mu^-\) is simply the annihilation of a quark and an antiquark into a virtual \(Z\) boson or photon, followed by the decay of the produced boson into a muon pair:

\[
q(p_1) + \bar{q}(p_2) \rightarrow \mu^-(k_1) + \mu^+(k_2)
\] (2.45)
with each particle’s four-momentum indicated as $p$ for initial-state particles, and $k$ for final-state particles.

At this point it is useful to introduce the Mandelstam variables:

$$\hat{s} \equiv -(p_1 + p_2)^2 = -(k_1 + k_2)^2$$  (2.46)
$$\hat{t} \equiv -(p_1 - k_1)^2 = -(k_2 - p_2)^2$$  (2.47)
$$\hat{u} \equiv -(p_1 - k_2)^2 = -(k_1 - p_2)^2$$  (2.48)

Of the three, $\hat{s}$ is the most intuitive as it corresponds to the centre-of-mass energy of the $q\bar{q}$ collision. At LO, this also corresponds to the square of the invariant mass $m_{\mu^+\mu^-}$ of the muon pair; in other words, $m_{\mu^+\mu^-}^2 = |\hat{s}|$ where $r$ is the four-momentum of the gauge boson.

The relevant parts of the interaction Lagrangian $\mathcal{L}_{Z/\gamma\psi}$ from Equation 2.28 are

$$\mathcal{L}_{Z/\gamma\mu} = ig EM \gamma^\rho Q_{\mu} q A_\rho + ig Z q \gamma^\rho \left(\frac{1 + \gamma^5}{2} g_{\mu L} + \frac{1 - \gamma^5}{2} g_{\mu R}\right) \mu^- Z_{\rho}$$  (2.49)

where the space-time index $\rho$ is used instead of the usual $\mu$ to avoid confusion with the particle.

The matrix element associated with the diagram in Figure 2.3 can now be calculated using the Feynman rules, which convert each external line, internal line and vertex into a mathematical expression, and give a prescription for combining these factors to form the matrix element. The complete list of Feynman rules can be found in references such as Ref. [58]. The result is:

$$\mathcal{M} = \mathcal{M}_\gamma + \mathcal{M}_Z$$  (2.51)

$$\mathcal{M} = [\mu^+(k_2) \gamma^\rho g EM Q_{\mu} \mu^-(k_1)] [\bar{q}(p_2) \gamma^\nu g EM Q_{\bar{q}} q(p_1)] \frac{\eta_{\rho\nu}}{r^2 - i\epsilon}$$
$$+ [\mu^+(k_2) \gamma^\rho g_{\mu} z \gamma^\nu g_{\mu L} + \frac{1 - \gamma^5}{2} g_{\mu R}] \mu^- (k_1)$$
$$\times [\bar{q}(p_2) \gamma^\rho g_{\nu} z \gamma^\nu g_{\nu L} + \frac{1 - \gamma^5}{2} g_{\nu R}] q(p_1)$$
$$\times \frac{1}{r^2 + M_Z^2 - iM_Z \Gamma_Z} \left(\eta_{\rho\nu} + \frac{r_{\rho} r_{\nu}}{M_Z^2}\right)$$  (2.52)

To enter the expression for the cross-section, the square of this matrix element must then be averaged over the four possible initial spin states and summed over the final spin states. The general result is simplified in the ultra-relativistic limit, where $\hat{s} \gg M_Z^2$ for all $\tilde{\psi}$ except the top quark: then all
masses in this calculation except $M_Z$ can be neglected. After using the relations

$$- \sum_{\text{spins}} \psi(p) \bar{\psi}(p) = i \gamma^\rho p_\rho - M_\psi \rightarrow i \gamma^\rho p_\rho$$  \hspace{1cm} (2.53)$$

$$- \sum_{\text{spins}} \bar{\psi}(k) \psi(k) = i \gamma^\rho k_\rho + M_\psi \rightarrow i \gamma^\rho k_\rho$$  \hspace{1cm} (2.54)$$

and taking all the traces involving the Dirac matrices, the result is

$$|\mathcal{M}|^2 = |A_{LL}|^2 \hat{s}^2 + |A_{RR}|^2 \hat{t}^2 + |A_{LR}|^2 \hat{u}^2 + |A_{RL}|^2 \hat{t}^2$$  \hspace{1cm} (2.55)$$

where

$$A_{ij} = g_{\text{EM}}^2 \frac{Q_i Q_j \hat{s}}{\hat{s}} + g_Z^2 \frac{g_\psi g_{\mu j}}{\hat{s}} - M_Z^2 - i M_Z \Gamma_Z$$  \hspace{1cm} (2.56)$$

Introducing this expression into Equation 2.44 and integrating yields the cross section

$$\sigma(q\bar{q} \rightarrow Z/\gamma' \rightarrow \mu^+\mu^-) = \frac{\hat{s}}{48\pi} \left( |A_{LL}|^2 + |A_{RR}|^2 + |A_{LR}|^2 + |A_{RL}|^2 \right)$$  \hspace{1cm} (2.57)$$

which is inversely proportional to $\hat{s}$ for $\hat{s} \gg M_Z^2$, since then $|A_{ij}|^2 \propto 1/\hat{s}^2$. The interference between the individual contributions to the cross section due to the photon and the $Z$ boson is manifest, as cross terms appear when squaring the matrix element. Similar interference effects can be important in calculations involving new spin-1 particles such as $Z'$ bosons discussed in Section 3.2.

One more step is necessary for the result of this calculation to become useful experimentally: after all, the initial state of collisions at the LHC consists not simply of quarks, but of protons. In addition to their three “valence quarks”, two up quarks and one down quark, protons are also made of gauge bosons (mainly gluons) binding these valence quarks together, as well as quark-antiquark pairs which are continually created and annihilated in interactions with these internal gauge bosons. This situation is illustrated in Figure 2.5. The quarks, antiquarks and gauge bosons making hadrons like the proton are collectively called “partons”.

Figure 2.5: Schematic illustration of parton interactions within a proton. Figure credit: DESY.
It follows that in a typical proton-proton collision, one parton from each incoming proton interact in the main process of the event. The probabilities for the identity of each parton are parametrized by Parton Distribution Functions (PDFs), which depend on the momentum transfer of the collision and the fraction \( x \) of the proton momentum carried by each type of parton. They are evaluated using data from deep inelastic scattering experiments and from hadron colliders. For example, the PDFs of the proton as evaluated by the MSTW collaboration \[132\] for \( |r^2| = 10 \text{ GeV}^2 \) and \((100 \text{ GeV})^2\) are shown in Figure 2.6.

**Figure 2.6:** Proton NLO PDFs from the MSTW collaboration at two momentum transfer scales: \(|r^2| = 10 \text{ GeV}^2\) and \((100 \text{ GeV})^2\) \[132\].

The LO cross section for the process \( pp \to Z/\gamma^* \to \mu^+\mu^- \) must then include the sum over all possibilities for the incoming quarks, weighted by their PDF. The Mandelstam variable \( s \) is now taken to represent the centre-of-mass energy of the proton-proton collision, and the quantity that must enter the cross section calculation is therefore \( x_1x_2s \), where \( x_1 \) and \( x_2 \) represent the fraction of the proton momentum carried by the incoming partons. In addition, the average over quark colours must be taken, because the quark-antiquark annihilation can only occur if they have the same colour. This results in an additional factor of \( 1/3 \). With these changes,

\[
\sigma(q\bar{q} \to Z/\gamma^* \to \mu^+\mu^-) = \frac{x_1x_2s}{144\pi} \left(|A_{LL}|^2 + |A_{RR}|^2 + |A_{LR}|^2 + |A_{RL}|^2\right) \tag{2.58}
\]

where the change \( \hat{s} \to x_1x_2s \) also modifies the \( A_{ij} \) from Equation 2.56.
One last point to consider is that because in general $x_1 \neq x_2$, the mediating boson will not be at rest in the centre-of-mass rest frame of the proton-proton system, but rather move along the beam axis $z$ with rapidity

$$y = \frac{1}{2} \ln \left( \frac{E_r + r_z}{E_r - r_z} \right) = \frac{1}{2} \ln \frac{x_1}{x_2}$$  \hspace{1cm} (2.59)

Finally, integrating over the PDF $f_q(x, r^2)$ of both protons while imposing four-momentum conservation gives the following expression for the differential cross section as a function of the momentum transfer and rapidity:

$$\frac{d\sigma}{d^2 r d y}(pp \rightarrow Z/\gamma^* \rightarrow \mu^+ \mu^-) = \int_0^1 dx_1 dx_2 \sum_q (f_{\bar{q}}(x_1, r^2) f_q(x_2, r^2) + f_q(x_1, r^2) f_{\bar{q}}(x_2, r^2))
\times \frac{1}{144\pi} (|A_{LL}|^2 + |A_{RR}|^2 + |A_{LR}|^2 + |A_{RL}|^2) \delta \left( |r^2| - x_1 x_2 s \right) \delta \left( y - \frac{1}{2} \ln \frac{x_1}{x_2} \right)$$  \hspace{1cm} (2.60)

where as above $A_{ij} = A_{ij}(r^2) \propto 1/|r^2|$ for $|r^2| \gg M_Z^2$.

This is an important result. It describes at Leading Order the behaviour of the Drell-Yan spectrum above the $Z$ peak in $m_{\mu^+\mu^-}$, where this process constitutes the dominant background to the search for new neutral resonances decaying into muon pairs.

**State-of-the-art Standard Model calculations and comparison to measurements**

In practice, state-of-the-art cross section calculations at higher orders are not performed analytically, but rather use Monte Carlo (MC) simulation techniques including contributions from higher orders in perturbative quantum field theory. This technique is described in more detail in Section 6.2.1.

Figure 2.7 shows the Standard Model cross section predictions for selected outcomes of proton-antiproton and proton-proton collisions. With the exception of the total cross section $\sigma_{\text{tot}}$, which is based on a parametrization from Ref. [52], they are calculated at NNLO in perturbative QCD with NLO electroweak corrections applied.

The Standard Model is in fact remarkably accurate in predicting the rates of production and decay of particles in high-energy collisions. Key examples of this are illustrated in Figure 2.8, which shows the outcome of several Standard Model production cross section measurements performed by the ATLAS experiment, compared with the corresponding theoretical expectations.
Figure 2.7: Standard Model cross section predictions for proton-antiproton (shown for $\sqrt{s} < 4$ TeV) and proton-proton (shown for $\sqrt{s} \geq 4$ TeV) collisions, at NNLO in perturbative QCD with NLO electroweak corrections applied. Vertical lines indicate the centre-of-mass energy at the Tevatron ($\sqrt{s} = 1.96$ TeV), and at the LHC ($\sqrt{s} = 7, 8$ and $13$ TeV). Figure credit: W. J. Stirling.
Figure 2.8: Summary of several Standard Model production cross section measurements performed by the ATLAS collaboration, compared with the corresponding theoretical expectations calculated at NLO or higher [33].


2.3 Limitations of the Standard Model

In spite of the outstanding success of the Standard Model of particle physics, there are strong experimental and theoretical indications that it is far from being a complete theory of physics at the most fundamental level.

2.3.1 Quantum Gravity

First, perhaps the most obvious shortcoming of the Standard Model is that it neglects gravitational effects. In Einstein’s general relativity, matter and force-carrying particles exist in curved spacetime, and this curvature is in turn affected by the mass-energy distribution. Quantum field theories like the Standard Model allow calculations for processes involving particles in curved spacetime [100], but not for the quantum fluctuations of fields to affect the curvature. Our current understanding of quantum gravitational effects would have such effects require non-renormalizable terms in the Lagrangian; in other words, quantum gravity is currently an effective field theory.

Like in the case of Fermi’s interaction, a non-renormalizable four-fermion contact interaction successful in describing radioactive decays, which were later understood as being mediated by W bosons, effective field theories are valid up to a certain energy scale, or equivalently down to a minimum distance at which effects beyond the ones described by the theory become important. For example, calculations using Fermi’s interaction are valid for energies $E \ll M_W$.

The presence of non-renormalizable terms in quantum theories of gravity might be indicative of an underlying, more complete theory, the details of which remain to be understood. Efforts are ongoing to design such a theory and reconcile general relativity with quantum field theory.

2.3.2 Dark Matter and Dark Energy

On the experimental side, astronomical observations indicate that ordinary matter as described by the Standard Model only represents a small fraction of the mass-energy inventory of the Universe. One of the first indications towards this has been seen in the rotation velocity distributions of stars in spiral galaxies. By looking at the visible distribution of matter in these galaxies, it is possible to calculate an expected velocity distribution for stars as a function of their distance from the galaxy centre. This expected velocity distribution does not match observations: stars outside the galactic bulge rotate faster than predicted [54, 149, 166], which hints towards the presence of invisible matter, named “dark matter”, that would be present in every galaxy.

The most definite evidence for dark matter comes from observations in the Bullet Cluster [74, 130], shown in Figure 2.9. This cluster is composed of two parts that have collided, and it is observed that the distribution of visible matter is manifestly different from the distribution of invisible matter inferred from gravitational lensing. In fact, the centre of the total mass is calculated to be offset from the centre of visible mass with a statistical significance of $8\sigma$ [75]. This
phenomenon is attributed to interactions between the visible matter constituents during the collision having slowed down these components of the cluster, while the weakly-interacting invisible matter passed unhindered.

Figure 2.9: The Bullet Cluster as seen by Magellan in the visible spectrum [75], with the distribution of x-ray emitting gas as seen by Chandra shown in red [129], and the distribution of invisible matter inferred from gravitational lensing shown in blue [75]. Figure credit: X-ray: NASA/CXC/CfA/M.Markevitch et al.; Lensing Map: NASA/STScI; ESO WFI; Magellan/U.Arizona/D.Clowe et al.; Optical: NASA/STScI; Magellan/U.Arizona/D.Clowe et al.

Another astronomical observation requiring an explanation beyond the Standard Model is that the expansion of the Universe is accelerating [142, 148], which points towards a very small but positive value for the cosmological constant of general relativity. This constant is generally understood as a form of energy, called “dark energy”, inherent to space itself, in other words the energy density of the vacuum.

Indications of the existence of dark matter and dark energy are also observed in the cosmic microwave background, as successively observed with increasing precision by the COBE, WMAP and Planck [143] satellites. The temperature map measured by these experiments is in fact very uniform across the sky, but temperature differences of $O(100 \, \mu K)$ exist between different directions. The distribution of these temperature fluctuations is sensitive to the composition of the Universe, and the Planck data are compatible with a dark matter content of 26.8% and a dark energy content of 68.3%, with ordinary matter representing only 4.9% of the mass-energy inventory of the Universe [144]: another clear indication that the Standard Model is not a complete description of nature.
2.3.3 Charge-Parity Violation and Matter-Antimatter Asymmetry

Again from the realm of astrophysics, a curious feature of the observable Universe is that most of the ordinary matter consists of particles, as opposed to antiparticles. This is surprising from the particle physics viewpoint if the initial state of the Universe is assumed to be a state of pure energy, because then matter and antimatter would have been produced in equal amounts after the Big Bang, implying that equal amounts of matter and antimatter should be observed.

One way to resolve this paradox involves interactions in particle physics that break the matter-antimatter symmetry, which implies violating the Charge-Parity (CP) symmetry, or equivalently time-reversal symmetry. There are indeed sources of such CP asymmetry in the flavour-mixing features of the Standard Model, but their strength is not large enough to explain the observed imbalance. Therefore, if the assumption about the initial state of the Universe holds, new physical processes, perhaps only occurring at very-high energies, must exist to contribute to this asymmetry.

Such new physical processes would correspond to additional terms in the Lagrangian. In the Standard Model, all the renormalizable interactions that could appear given the model’s particle content are effectively realized in nature, with only one exception: an additional gluon interaction term could have been present in $\mathcal{L}_{G\nu}$ in Equation 2.1

$$\mathcal{L}_{G\nu} = -\frac{g_3^2 \Theta_3}{64\pi^2} \epsilon_{\mu\nu\lambda\rho} G^{\mu\nu} G^{\alpha\lambda\rho}$$

(2.61)

Similar terms for the electroweak gauge bosons would have no physical effect [58], but this gluon interaction term would constitute an additional source of CP symmetry violation. However, the existence of this term would imply a non-zero value for the electric dipole moment of the neutron $\delta_n$, on which there are stringent experimental bounds. The best constraint to date is $|\delta_n| < 2.9 \times 10^{-26} e\cdot cm$ at 90% Confidence Level (CL) [40]. This implies that the $\Theta_3$ term is effectively absent: $|\Theta_3| \lesssim 10^{-10}$ [52]. The matter-antimatter asymmetry of the Universe therefore remains a mystery.

2.3.4 Neutrino Masses and Flavour Oscillation

While in the Standard Model neutrinos have exactly zero mass, experiments worldwide have demonstrated that neutrinos are in fact massive, following from the observation of neutrino flavour oscillations [111]. In fact, such oscillations are only possible if at least two neutrino flavours are massive. As a consequence, leptons in the mass basis are rotated from their representations in the flavour basis, just like in the quark sector. The charged leptons are taken to remain the same, and neutrinos in the mass basis are labelled $\nu_1$, $\nu_2$ and $\nu_3$. The equivalent of the CKM matrix is called the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix [52]. The off-diagonal elements of the PMNS matrix are large, implying that mixing is much more pronounced for neutrinos than for quarks.

A possible mechanism to generate neutrino masses is called the see-saw mechanism. The main
idea is that it is possible to add right-handed neutrinos to the Standard Model, transforming as singlets: \((1, 1, 0)\) under the Standard Model symmetries. This implies that right-handed neutrinos do not interact via any of the interactions of the Standard Model, and are therefore undetectable except indirectly because of their gravitational interactions. Combined with the fact that their masses could be large, this incidentally makes right-handed neutrinos a possible dark matter candidate.

With right-handed neutrinos \(N_m\), more terms appear in the Lagrangian, including a Majorana mass term. Such a term is allowed because right-handed neutrinos are singlets under all \(SU(3)_C \times SU(2)_L \times U(1)_Y\) interactions, and can therefore have an explicit mass without breaking gauge invariance. The Lagrangian terms involving right-handed neutrinos are the following:

\[
\mathcal{L}_N = -\frac{1}{2} \bar{N}_m (\gamma^\mu \partial_\mu + M_m) N_m - \frac{y_{N,mn}}{\sqrt{2}} (v + H) \bar{\nu}_m N_n
\]

(2.62)

To investigate the consequences of this Lagrangian, consider a simplified case where \(y_{N,mn}\) is diagonal. Then for each generation the neutrino mass matrix has the form

\[
\begin{bmatrix}
0 & \frac{1}{\sqrt{2}} y_{N,1} \\
\frac{1}{\sqrt{2}} y_{N,1} & M
\end{bmatrix}
\]

(2.63)

The positive-definite eigenvalues of this matrix are

\[
\frac{1}{2} \left| M \pm \sqrt{M^2 + 2y_{N}^2v^2} \right| \rightarrow \frac{y_{N}^2v^2}{2M}, M
\]

(2.64)

with the values on the right in the limit \(M \gg y_{N} v\).

The largest of these eigenvalues, \(M\), corresponds to the mass of right-handed neutrinos, and a very small mass is given to left-handed neutrinos in the process. Given the experimental upper bound on the sum of the masses of left-handed neutrinos of 0.23 eV at 95% CL [52], this means that the mass of right-handed neutrinos could range between \(O(100 \text{ GeV})\), for Yukawa couplings \(y_N \sim y_L \sim 10^{-6}\), up to \(O(10^{14} \text{ GeV})\) for \(y_N \sim 1\).

The experimental study of neutrino oscillations and direct measurement of their masses is a very rich field of study, however further discussion of this area is outside the scope of this dissertation. At any rate, the existence of non-zero masses for the left-handed neutrinos is a clear indication of physics beyond the Standard Model.

### 2.3.5 Hierarchy, Fine-Tuning and Elegance

The hierarchy problem is a theoretical problem, raising the question of the existence of such a wide range of energy scales in nature. At the lower end of the spectrum, neutrino masses are thought to be of \(O(0.1 \text{ eV})\), while the masses of other fermions range from the electron mass at 511 keV, to the top quark mass at 173 GeV. In the Standard Model, the latter range is understood as arising from
the differences in strength of the Yukawa couplings between the Higgs boson and fermions, but the origin of the precise values of these couplings is a mystery.

The electroweak symmetry breaking scale is $\mathcal{O}(100 \text{ GeV})$ like the top quark mass, but then the scale of gravity is the Planck scale, $\mathcal{O}(10^{19} \text{ GeV})$. This is equivalent to the observation that while the three forces present in the Standard Model have couplings that are only a little smaller than unity, gravity couples to matter in an extremely weak way, for reasons unknown.

Linked to this hierarchy problem is the Higgs fine-tuning problem. While the Higgs mass is a well-defined quantity in the Standard Model, this parameter receives important radiative corrections, linked to the fact that the Higgs boson is assumed to be a fundamental scalar particle. The largest of these corrections comes from top quark loops. In the absence of new, currently undiscovered particles to cancel these quantum corrections, the natural value for the mass of the Higgs boson would be up at the Planck scale, only to be brought back to the observed value using an ad hoc correction spanning 16 orders of magnitude. It is therefore theorized that the quantum corrections to the Higgs mass must be cancelled by contributions from new particles, most naturally with masses near the electroweak symmetry breaking scale. Other, more radical solutions to this problem also exist, and are described in Section 3.1.

Finally, the Standard Model has 19 degrees of freedom, a number that grows to 26 if neutrino masses and mixing angles are included. It is commonly thought that the simplest most fundamental theory of nature should have a smaller number of free parameters: this aesthetic elegance requirement is yet another motivation to look for physics beyond the Standard Model.
Chapter 3

Beyond the Standard Model

Many hypotheses reach beyond the Standard Model in attempts to address the limitations discussed in Section 2.3. This chapter first presents an overview of such hypotheses that predict new high-mass dilepton resonances in Section 3.1. Particular attention is devoted to models postulating new gauge symmetries in Section 3.2.

3.1 Hypotheses Beyond the Standard Model Predicting New High-Mass Resonances

Searches for new high-mass bosons decaying into lepton pairs are motivated by a wide variety of new models, each trying to address a subset of the Standard Model’s limitations. Many of these models postulate that the gauge group of the Standard Model does not represent the most fundamental symmetry realized in nature, but rather that the forces in the Standard Model are unified at high energies under a larger gauge symmetry. The possibility of new gauge symmetries is explored in more detail in Section 3.2.

Alternatively, many models hypothesize the existence of additional dimensions of space-time: allowing gravity to propagate in these extra dimensions then explains why gravity is perceived as much weaker than the other forces of nature. As a result, the true strength of gravity is comparable to that of other forces, the actual Planck mass is comparable to the electroweak symmetry breaking scale and there is no hierarchy problem. Popular examples of this class of models include the Randall-Sundrum models [147], which predict excited states of the graviton that can decay into lepton pairs, and the models proposed by Arkani-Hamed, Dimopoulos and Dvali (ADD) [9]. If the energy available at the collider is sufficiently above the actual Planck mass, quantum black holes may also be produced, and decay in the dilepton final state [94, 106]. Variants of these extra-dimensional models, as well as TeV$^{-1}$ models [7, 8, 39, 49], also allow gauge bosons to propagate in the additional dimensions, predicting excited gauge boson states. Of these, the Kaluza-Klein spin-1 bosons $Z_{KK}$ and $\gamma_{KK}$ may also decay into lepton pairs.
Technicolour models propose another solution to the Higgs fine-tuning problem: to break electroweak symmetry using new strong dynamics. Many technicolour models are in fact incompatible with the properties of the newly-discovered Higgs boson, but some, such as Minimal Walking Technicolour, include a composite Higgs boson with the right mass, spin and parity [93, 98, 99]. Since the Higgs boson is then a composite state, there is no fundamental scalar particle. This model also predicts other bound states called technimesons, which may be detected as resonances in the dilepton spectrum.

Additional models predict new particles decaying into lepton pairs: models with new chiral bosons $W^*$ and $Z^*$ [68–71], torsion models [50, 91, 145, 155], $R$-parity violating supersymmetry [43], and many others, further motivating the search. There is also the possibility of new particles that no one has yet imagined!

In this dissertation, the interpretation of the search for new high-mass resonances decaying into muon pairs will concentrate on models with new gauge symmetries; these are described in more detail in the next section.

3.2 New Gauge Symmetries

3.2.1 The Sequential Standard Model

Arguably the most simple extension to the Standard Model involves force-carrying particles with the same couplings to fermions, the only difference between the new particles and the existing ones being their predicted higher mass. Such a model is called the Sequential Standard Model (SSM) [123], and predicts new spin-1 bosons $W^\prime_{\text{SSM}}$ and $Z^\prime_{\text{SSM}}$, directly analogous to the $W$ and $Z$ bosons from the Standard Model. While the SSM is poorly motivated from a theoretical perspective, it provides a useful benchmark to compare the relative sensitivity of searches in different channels or from different experiments.

Unlike the $W$ and $Z$ bosons, which interact with each other in cubic and quartic interaction terms, the couplings of the bosons introduced in the SSM with each other and with the $W$ and $Z$ are assumed to be zero. This arbitrary choice is made for all $Z^\prime$ bosons under study, in order to avoid introducing additional model-dependent degrees of freedom into the interpretation of the search.

3.2.2 Models from Grand Unified Theories

More interesting than the SSM is the possibility of new vector bosons from Grand Unified Theories (GUT). In GUT, the gauge group of the Standard Model is extended, usually to a simple group, to provide a unifying picture of gauge interactions in which all fermions belong to the same multiplet representation of the group. These models are motivated by the fact that gauge couplings in the Standard Model seem to converge towards the same value at very high energy scales, as illustrated
in Figure 2.2. Further, GUT typically feature a reduction in the number of free parameters in the theory, as well as natural neutrino masses due to the presence of right-handed neutrinos, and sometimes a cancellation of quadratic divergences of the Higgs mass in an attempt to solve the fine-tuning problem.

Perhaps one of the most popular examples of GUT in high-energy particle physics is the superstring-motivated $E_6$ model [127]. In particular, in a superstring theory with 10 space-time dimensions and an $E_8 \times E_8'$ internal gauge symmetry, compactification to four dimensions reduces one of the $E_8$ to $E_6$. This $E_6$ comprises the gauge symmetries of the Standard Model, while the $E_8'$ interacts only gravitationally with ordinary matter and forces. While this hypothesis is fairly speculative and may not correspond to what actually occurs in nature, studying it is of particular interest because the predictions originating from its many different possible symmetry-breaking scenarios do cover many of the different phenomenological possibilities explored in the literature. For instance, alternative symmetry-breaking scenarios such as $E_8$ first breaking into $SO(16)$ instead [45] give similar predictions at the energy scale explored by the LHC.

The breakdown of the $E_6$ gauge symmetry to $SO(10) \times U(1)$ is particularly appealing because then for each of three generations of matter, the Standard Model fermions plus a right-handed neutrino are part of the same representation of $SO(10)$, while eleven additional exotic fermions form a $10$ and a singlet. These exotic fermions can be given masses as high as the symmetry-breaking scale of $E_6$, and will not be discussed further. While it is also possible for the $E_6$ symmetry to break down following a different path, explicitly either $E_6 \rightarrow SU(6) \times SU(2)$ or $E_6 \rightarrow [SU(3)]^3$, the main experimental prediction that follows from these scenarios is qualitatively the same: a $Z'$ boson may exist at the TeV scale. All possibilities are explored in detail by London and Rosner in Ref. [127].

$SO(10)$ can then break down to the Standard Model following one of two scenarios: the Georgi-Glashow Model or the Pati-Salam Model.

### The Georgi-Glashow Model

A first symmetry-breaking scenario is via $SU(5)$:

$$E_6 \rightarrow SO(10) \times U(1)_\psi$$

$$\rightarrow SU(5) \times U(1)_X \times U(1)_\psi$$

$$\rightarrow SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_\theta$$

$$\rightarrow SU(3)_C \times SU(2)_L \times U(1)_Y$$

$$\rightarrow SU(3)_C \times U(1)_Q$$ (3.1)

Here all the symmetries of the Standard Model become unified at higher energies into the simple group $SU(5)$. In the $SU(5)$ model, first explored by Georgi and Glashow [105], the representations
containing the fermions from the Standard Model are one $\bar{5}$ and one $10$ for each generation:

$$
\begin{bmatrix}
\bar{d}_s \\
\bar{d}_g \\
\bar{d}_b \\
e^-
\end{bmatrix}
\begin{bmatrix}
0 & \bar{u}_b & -\bar{g}_d & -d_f \\
-\bar{u}_b & 0 & \bar{u}_r & -u_g \\
\bar{g}_d & -\bar{u}_r & 0 & -u_b & -d_b \\
u_r & u_g & u_b & 0 & -e^+
\end{bmatrix}
\begin{bmatrix}
d_r \\
d_g \\
d_b \\
e^+
\end{bmatrix}_L
$$

(3.2)

The gauge bosons from the adjoint representation $24$ of $SU(5)$ include the Standard Model gauge bosons, as well as leptoquarks. The latter can mediate experimentally unseen processes such as proton decay (e.g. via $p \to \pi^0 e^+$) and $K^0 \to \mu^\pm e^\mp$. For this reason, the $SU(5)$ breaking scale has to occur at very high energies, thereby giving large masses to the leptoquarks.

Nevertheless, the spontaneous symmetry breaking of $E_6$ to $SU(5)$ results in two additional $U(1)$ gauge groups named $U(1)_\psi$ and $U(1)_\chi$, and the lightest linear combination of the associated massive gauge bosons $Z'_\psi$ and $Z'_\chi$ could be accessible at the LHC if the symmetry breaking of $SU(5)$ occurs in two steps, with the breaking of $U(1)_\theta$ involving a vacuum expectation value at the TeV scale. The experimental signature would then be a new gauge boson $Z'_\theta$, where

$$
Z'_\theta = Z'_\psi \cos \theta_{E_6} + Z'_\chi \sin \theta_{E_6}.
$$

(3.3)

The couplings of $Z'_\theta$ to fermions are detailed in Table I of Ref. [127].

**The Pati-Salam Model**

An even more interesting symmetry-breaking scenario involves the Pati-Salam Model [141], which in turn breaks down to the Standard Model via the minimal left-right symmetric model as follows:

$$
E_6 \rightarrow SO(10) \times U(1)_\psi \\
\quad \rightarrow SU(4)_C \times SU(2)_L \times SU(2)_R \\
\quad \rightarrow SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \\
\quad \rightarrow SU(3)_C \times SU(2)_L \times U(1)_Y \\
\quad \rightarrow SU(3)_C \times U(1)_Q
$$

(3.4)

This particular mechanism makes a clear separation between the flavour and colour sectors of the theory. On the flavour side, two triplets of weak bosons are present: the $W$ and $Z$ bosons from the Standard Model $SU(2)_L$, but also $W'_R$ and $Z'_R$ from $SU(2)_R$, making the left-right symmetry explicit. On the colour side, the strong force symmetry group is $SU(4)_C$, with leptons joining quarks in colour quadruplets (charged leptons with down-type quarks, and neutrinos with up-type quarks). This unification also predicts leptoquarks, which gain large masses when $SU(4)_C$ is broken into the
familiar $SU(3)_C$ and the Abelian $U(1)_{B-L}$ of the minimal left-right symmetric model [46, 136, 154] where $B$ is the baryon number, $L$ is the lepton number and $B - L$ is the conserved quantum number.

The Standard Model is recovered when the symmetry $SU(2)_R \times U(1)_{B-L}$ breaks down to $U(1)_Y$, in a manner very similar to the electroweak symmetry breaking $SU(2)_L \times U(1)_Y \rightarrow U(1)_Q$. If this last step of the breakdown to the Standard Model happens near the TeV scale, then the $W'_R$, $Z'_R$ and/or $Z'_{B-L}$ bosons could be observable at the LHC.

### Minimal $Z'$ Models

The Minimal $Z'$ Models [151] constitute a parametrization aiming to describe a large section of the $Z'$ phenomenology potentially accessible at the LHC using only three parameters beyond these of the Standard Model. The new parameters are the $Z'_{\text{Min}}$ mass $M_{Z'_{\text{Min}}}$, a coupling strength $\gamma'$ and a mixing angle $\theta_{\text{Min}}$. The latter two are related to the $Z'_{\text{Min}}$ coupling $g_Y$ to the electroweak hypercharge $Y$, and its coupling $g_{B-L}$ to the $B - L$ current, as follows:

\begin{align}
\frac{g_Y}{g_Z} &= \frac{\gamma'}{\sin \theta_{\text{Min}}} \\
\frac{g_{B-L}}{g_Z} &= \frac{\gamma'}{\cos \theta_{\text{Min}}}
\end{align}

where $g_Z = 2M_Z/v$, and $v = 246$ GeV is the Higgs vacuum expectation value in the Standard Model. $\gamma'$ therefore measures the overall coupling strength of the $Z'_{\text{Min}}$ boson relative that of the $Z$ boson.

As in the case of $\theta_{E_6}$, certain values of $\theta_{\text{Min}}$ correspond to specific $Z'$ bosons: in this case they are $Z'_{B-L}$, $Z'_R$ and $Z'^\chi$. The values of $\gamma'$ and $\theta_{\text{Min}}$ for these three specific models are shown in Table 3.1. The $Z'^\chi$ boson corresponds to the point in parameter space shared between Minimal $Z'$ Models and the $\theta_{E_6}$ parametrization.

**Table 3.1:** Values of $\gamma'$ and $\theta_{\text{Min}}$ for three specific models: $Z'_{B-L}$, $Z'_\chi$ and $Z'_R$.

<table>
<thead>
<tr>
<th>$\gamma'$</th>
<th>$Z'_{B-L}$</th>
<th>$Z'_\chi$</th>
<th>$Z'_R$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\frac{\sqrt{5}}{8} \sin \theta_W$</td>
<td>$\frac{\sqrt{41}}{21} \sin \theta_W$</td>
<td>$\frac{\sqrt{25}}{12} \sin \theta_W$</td>
<td></td>
</tr>
<tr>
<td>$\sin \theta_{\text{Min}}$</td>
<td>0</td>
<td>$-\sqrt{\frac{16}{41}}$</td>
<td>$-\sqrt{\frac{4}{5}}$</td>
</tr>
<tr>
<td>$\cos \theta_{\text{Min}}$</td>
<td>1</td>
<td>$\sqrt{\frac{25}{41}}$</td>
<td>$\sqrt{\frac{1}{5}}$</td>
</tr>
</tbody>
</table>

### Summary

To summarize, in Grand Unified Theories all forces of nature originate from a unified gauge symmetry. In the process of breaking this symmetry, additional $SU(2)$ and $U(1)$ groups appear at scales above that of the Standard Model. All additional $U(1)$ symmetries would manifest themselves via the production of $Z'$ gauge bosons, while an additional $SU(2)$ symmetry would imply the exis-
tence of both $W'$ and $Z'$ bosons. These additional gauge bosons might be observable directly at high-energy colliders or indirectly in precision experiments.

### 3.2.3 Experimental Limits on New Gauge Bosons

Previous experimental limits at 95% CL on the mass of new gauge bosons $Z'$ are shown in Table 3.2. The earliest TeV-scale limits on the existence of $Z'$ bosons came indirectly via limits on four-fermion contact interactions [67]. Such interactions would interfere with fermion pair-production occurring via the Drell-Yan process, and impact the production cross sections and angular distributions of the final-state fermions. The strongest indirect limits are set by the experiments at the Large Electron-Positron Collider (LEP) [2, 92, 120, 140], which exclude at 95% CL a $Z_{SSM}'$ boson with $M_{Z'} < 1.787$ TeV [124].

In their time, the CDF and D0 experiments at the Tevatron have set the strongest direct limits on the existence of new gauge bosons, using proton-antiproton collision data at $\sqrt{s} = 1.96$ TeV. These limits are also shown in Table 3.2. The limits set by the D0 collaboration made use of 5.3 fb$^{-1}$ of data in the dielectron channel [90], and the limits set by the CDF collaboration made use of 4.6 fb$^{-1}$ of data in the dimuon channel [66].

Finally, because the Large Hadron Collider currently holds the record for the highest centre-of-mass energy attained in a particle collider, it comes as no surprise that the most stringent direct limits to date on the existence of new gauge bosons come from the ATLAS and CMS experiments. Using proton-proton collision data amounting to 5 fb$^{-1}$ at $\sqrt{s} = 7$ TeV and 4 fb$^{-1}$ at $\sqrt{s} = 8$ TeV [78], the latest published results by CMS exclude at 95% CL a $Z_{SSM}'$ boson with $M_{Z'} < 2.59$ TeV and a $Z_{\psi}'$ with $M_{Z'} < 2.27$ TeV. The most recent preliminary results from CMS using 20 fb$^{-1}$ of data collected at $\sqrt{s} = 8$ TeV [77] exclude at 95% CL a $Z_{SSM}'$ boson with $M_{Z'} < 2.96$ TeV and a $Z_{\psi}'$ with $M_{Z'} < 2.60$ TeV. Results obtained by the ATLAS experiment are discussed in detail in Chapter 8.

**Table 3.2:** Previous experimental limits at 95% CL on the mass of new gauge bosons $Z'$.

CMS does not provide explicit limits on $Z_{\chi}'$; they would be slightly better than the limits on $Z_{\psi}'$. Results by the ATLAS collaboration are not shown here, but in Table 8.1.

<table>
<thead>
<tr>
<th>Collaboration</th>
<th>Channel</th>
<th>$Z_{SSM}'$</th>
<th>$Z_{\psi}'$</th>
<th>$Z_{\chi}'$</th>
</tr>
</thead>
<tbody>
<tr>
<td>LEP combination</td>
<td>$e^+e^- \rightarrow e^+e^-, \mu^+\mu^-, \tau^+\tau^-, q\bar{q}$</td>
<td>1.787</td>
<td>0.481</td>
<td>0.673</td>
</tr>
<tr>
<td>D0</td>
<td>$p\bar{p} \rightarrow e^+e^-$</td>
<td>1.023</td>
<td>0.891</td>
<td>0.903</td>
</tr>
<tr>
<td>CDF</td>
<td>$p\bar{p} \rightarrow \mu^+\mu^-$</td>
<td>1.071</td>
<td>0.917</td>
<td>0.930</td>
</tr>
<tr>
<td>CMS</td>
<td>$pp \rightarrow e^+e^-, \mu^+\mu^-$</td>
<td>2.59</td>
<td>2.27</td>
<td>-</td>
</tr>
<tr>
<td>CMS (preliminary)</td>
<td>$pp \rightarrow e^+e^-, \mu^+\mu^-$</td>
<td>2.96</td>
<td>2.60</td>
<td>-</td>
</tr>
</tbody>
</table>
Chapter 4

The ATLAS Experiment at the Large Hadron Collider

A truly international endeavour, the Large Hadron Collider (LHC) [51, 56, 57, 97] is the largest and most powerful particle accelerator ever built. This chapter first describes in Section 4.1 the LHC accelerator chain within the CERN accelerator complex, which accelerates protons to achieve collisions with record centre-of-mass energy and very high rates at four interaction points.

Immense detectors are present at each of these interaction points to record the outcomes of these collisions and study the underlying physical processes: they are the ALICE, ATLAS, CMS and LHCb experiments. This dissertation makes use of data collected by the ATLAS experiment [12, 13], described in Sections 4.2 to 4.8

4.1 The Large Hadron Collider

The Large Hadron Collider is located at CERN near Geneva, Switzerland, in the 27-kilometre-long tunnel previously used by the Large Electron-Positron Collider (LEP) from 1989 to 2000. It currently holds the record for the highest energy achieved in an artificial particle accelerator, having successfully accelerated protons to an energy of 3.5 TeV in 2010 and 2011, and 4 TeV in 2012. Its design energy of 7 TeV per proton is projected to be attained in the next years. In comparison, the Tevatron accelerated protons and antiprotons to an energy of 0.98 TeV. Unlike the Tevatron, the LHC is a proton-proton accelerator; this choice has been made to maximize the intensity of the beam and therefore the instantaneous luminosity, related to the collision rate. The LHC is also used to accelerate heavy ions, such as lead nuclei.
4.1.1 Accelerator Chain

The accelerator chain leading to the main ring of the LHC consists of pre-existing machines, upgraded to meet the requirements of their new purpose as pre-accelerators. The chain starts at the proton source, a bottle of hydrogen gas. A duoplasmotron is used to ionize the hydrogen atoms, by bombarding them with electrons; electric fields then complete the separation, sending the resulting protons into the accelerator chain. The CERN accelerator complex is shown in Figure 4.1.

![Figure 4.1: The CERN accelerator complex. Figure credit: CERN.](image)

The protons are first accelerated to an energy of 50 MeV by a linear accelerator, LINAC 2, which consists of radiofrequency cavities containing an alternating electric field. The length of successive cavities increases, matching the desired speed of the protons, so that protons in phase with the electric field are uniformly accelerated, while others arriving later or earlier are accelerated respectively more or less. The net result is the formation of proton bunches, accelerating jointly. Quadrupole magnets ensure that the bunches remain focused in the transverse directions.
These proton bunches are then injected into the Proton Synchrotron Booster, the first of a series of synchrotrons, i.e. circular accelerators where the particles follow a fixed trajectory, with dipole magnets providing the magnetic field to curve the particles’ path. This magnetic field must increase as the particles are accelerated, in order to maintain the radius of the circular trajectory constant. The maximal energy attainable by proton synchrotrons is limited by the strength of their dipole magnets as compared with the radius of curvature imposed by the accelerator’s geometry. This implies that there are two ways to reach higher energies: stronger magnets and larger rings. The Proton Synchrotron Booster actually consists of four superimposed synchrotron rings which accelerate proton bunches to an energy of 1.4 GeV, before sending them into larger and more powerful synchrotrons.

The proton bunches then enter the Proton Synchrotron, which once was the world’s most powerful particle accelerator when it first started in 1959. With a circumference of 628 meters, it can accelerate protons up to an energy of 25 GeV. During its long service, this versatile machine has also served to accelerate heavier nuclei, as well as electrons, positrons and antiprotons.

The next accelerator in the chain is the Super Proton Synchrotron, with a circumference of nearly 7 km, bringing the beam energy to 450 GeV per proton. It is from proton-antiproton collision events obtained using this accelerator that the UA1 and UA2 experiments discovered the W and Z bosons in 1983 [161–164]. Finally, the proton bunches are injected into the LHC, in opposite directions, for their final acceleration and eventual collision.

4.1.2 Main Ring

The LHC makes use of the LEP tunnel, shown in Figure 4.2, which only has room for a single ring of magnets. But unlike LEP or the Tevatron which accelerated beams with particles of equal and opposite charge, allowing both beams to share the same ultra-high vacuum tube and magnet system, the LHC’s two proton beams require dipole magnets of opposite polarity.

The answer to this challenge is the “two-in-one” superconducting electromagnet design, illustrated in Figure 4.3. In this economical configuration, the magnets share the same envelope, iron yoke and services. The magnets, made of a niobium-titanium alloy, are cooled down to a temperature of 1.9 K using liquid helium. At this temperature, helium is in the superfluid state, which implies excellent heat conduction properties. The use of superconducting technology is necessary in order to reach the very high magnetic fields required: given the curvature imposed by the LHC layout, at the design energy of 7 TeV per proton the dipole magnets need to reach 8.33 T.

In addition to the dipole magnets responsible for bending the beam’s trajectory, quadrupole and sextupole magnets focus the beam in the transverse directions, while the longitudinal stability of proton bunches is provided by the superconducting radio-frequency cavities responsible for ac-

---

1 This is in contrast with a cyclotron such as the one at TRIUMF, where the magnetic field in the accelerator is constant but the radius of the particles’ trajectory increases as they are accelerated, resulting in an outward spiral path.
celerating the beam and compensating for any energy losses once the collision energy is attained. Additional magnets are used to improve the beam quality by compensating for second-order effects within the beams. Finally, at the interaction points, dedicated dipole magnets initiate beam recombination and ensure beam separation on the other side of the crossing point, while dedicated triplets of quadrupole magnets are responsible for focusing the beam to achieve the highest possible collision rate. The placement of these magnets around the ATLAS interaction point is shown in Figure 4.4.

![Magnets in the LHC tunnel](image1.jpg)

**Figure 4.2:** Magnets in the LHC tunnel. Figure credit: CERN.

![LHC magnet components](image2.jpg)

**Figure 4.3:** LHC magnet components. Figure credit: CERN.
Figure 4.4: Placement of the dipole and quadrupole magnets responsible for bringing the LHC beams into collision at the ATLAS interaction point. Also shown are the positions of the LUCID, Zero Degree Calorimeter and ALFA detectors discussed in Section 4.7 [12].

4.1.3 Delivered Luminosity and Beam Conditions in Run-I

Following the first proton-proton collisions obtained at a centre-of-mass energy of 900 GeV in November 2009, the accelerating power of the LHC was used to achieve a record energy of 2.36 TeV one week later. A centre-of-mass energy of 7 TeV was reached in March 2010, albeit with a low luminosity. In total, 48.1 pb\(^{-1}\) of proton-proton collisions were delivered to the ATLAS experiment that year, making possible the “re-discovery” of Standard Model processes and the first iteration of many searches for new physics at the LHC.

The luminosity delivered to ATLAS increased dramatically in 2011, reaching 5.46 fb\(^{-1}\) of collisions at \(\sqrt{s} = 7\) TeV by the end of the year. Then the centre-of-mass energy was raised to \(\sqrt{s} = 8\) TeV in 2012, when 22.8 fb\(^{-1}\) of collisions were delivered to the ATLAS detector. In total, \(1.80 \times 10^{15}\) proton-proton collisions were recorded by the ATLAS detector in Run-I of the LHC. The cumulative progression of the luminosity delivered to ATLAS is shown in Figure 4.5, along with the fraction of these data successfully recorded by ATLAS and passing all quality requirements, as discussed in Sections 4.7 and 4.8.

In fact, the highest instantaneous luminosity achieved by the LHC in 2012 was \(7.73 \times 10^{33}\) cm\(^{-2}\)s\(^{-1}\), close to the design value of \(10^{34}\) cm\(^{-2}\)s\(^{-1}\) = 10 nb\(^{-1}\)s\(^{-1}\). This was achieved with half the proton bunch crossing rate originally planned: collisions occurred every 50 ns, instead of the nominal time separation of 25 ns between bunches. This was compensated by a higher number of protons in each bunch: specifically, instead of accelerating 2808 bunches with \(1.15 \times 10^{11}\) protons per bunch, the LHC accelerated 1380 bunches with \(1.7 \times 10^{11}\) protons each.

This implies that the reconstruction of detected collision events benefited from a greater separation between bunch crossings, i.e. lower “out-of-time pileup”, but that a larger number of simultaneous proton-proton interactions, or “in-time pileup”, occurred at each bunch crossing. The hardware and event reconstruction software of ATLAS met this challenge, and operated successfully under these conditions. The distributions of the observed number of interactions per crossing in 2011 and 2012 are shown in Figure 4.6. Figure 4.7 displays a \(Z \rightarrow \mu^+ \mu^-\) candidate event observed by the ATLAS detector with reconstructed vertices from 25 simultaneous proton-proton interactions, a number in the bulk of the distribution for collisions recorded in 2012.
**Figure 4.5:** Cumulative luminosity versus time delivered to ATLAS (green), recorded by ATLAS (yellow), and certified to be good quality data (blue) during stable beams and for proton-proton collisions at 7 and 8 TeV centre-of-mass energy in 2011 and 2012 [33].

**Figure 4.6:** Luminosity-weighted distribution of the mean number of interactions per crossing for the dataset collected in 2011 and 2012 [33].
Figure 4.7: $Z \rightarrow \mu^+\mu^-$ candidate event with reconstructed vertices from 25 simultaneous proton-proton interactions in the ATLAS detector. The two muons, which come from the same vertex, are highlighted in yellow [33].
4.2 The ATLAS Detector: Overview

With its extraordinary collision rates and beam energy, the LHC offers a tremendous technological challenge for detector designers. The ATLAS detector, illustrated in Figure 4.8, was created to meet this challenge: it serves as one of the two general-purpose detectors at the LHC, with the ambitious goal to search for the Higgs boson from the Standard Model in addition to being sensitive to a very wide range of new physics that might appear at the TeV scale. It also provides an opportunity to investigate much of the Standard Model’s high-energy spectrum, studying processes involving the top quark and flavour physics as well as performing precision tests of QCD and electroweak interactions.

To achieve these goals, precise and fast measurements of the physical objects in the final state of collisions are essential, with the largest possible acceptance and efficiency for detecting signal processes. The next sections describe the technological features making such performance possible.

4.2.1 Detector Design

ATLAS is composed of four major subsystems, built in cylindrical layers: from the innermost to the outermost part, they are the Inner Detector (ID), the Electromagnetic Calorimeter, the Hadronic Calorimeters and the Muon Spectrometer (MS). Each of these subsystems is divided into three regions: the central region is called the “Barrel”, bordered by “Endcap” regions on each side. As illustrated in Figure 4.9, when particles produced in the LHC collisions interact with the detector, they produce distinctive patterns which make their identification possible.

The ID, described in Section 4.3, measures tracks from charged particles: the momentum of these tracks is determined by their curvature in the magnetic field provided by a solenoid magnet. As well, the information from these tracks makes possible the reconstruction of primary vertices, allowing to distinguish different collision points, and the reconstruction of secondary vertices coming for example from $B$-meson decays. The Electromagnetic Calorimeter, responsible for identifying and measuring the energy of electrons and photons, and the Hadronic Calorimeters, which measure the energy of hadrons while ensuring to stop them before they reach the MS, are briefly described in Section 4.4.

Particular attention is devoted to the MS described in Section 4.5, as it is the main subsystem used to detect and measure the momentum of muons, especially at very high momentum. One of the prominent features of ATLAS is its toroid magnet system, providing the strong magnetic field necessary to perform this measurement: it is also described there.

Section 4.6 defines the missing transverse momentum, a very useful observable making use of the information from all subsystems allowing the indirect detection of particles which do not interact with any of the detector components, such as neutrinos. Section 4.7 outlines the strategies used to monitor the luminosity delivered to ATLAS. Finally, sophisticated trigger and Data Acquisition
(DAQ) systems are responsible for the first stages of event selection, recording the outcome of experimentally interesting collision events and discarding the vast majority of the rest. As well, data quality monitoring ensures that all subsystems are working properly and that the collected data are suitable for analysis. These systems are described in Section 4.8.

### 4.2.2 Coordinate System

The ATLAS collaboration employs a right-handed coordinate system, with the \( z \)-axis along the beam direction, also called the longitudinal direction. The origin is defined as the nominal interaction point in the centre of the detector. The \( x \)-axis points from this origin toward the centre of the LHC ring, while the \( y \)-axis points upward. Cylindrical coordinates are often used, \( r \) being the radius in the transverse plane \((x, y)\) and \( \phi \) being the azimuthal angle around the beam pipe defined from the positive \( x \)-axis. When using spherical coordinates, the polar angle \( \theta \) is defined from the positive \( y \)-axis. The transverse momentum of particles is defined in terms of \( \theta \) and the momentum \( p \) as \( p_T = p \sin \theta \), and the pseudorapidity is defined as \( \eta = -\ln \tan(\theta/2) \).

When describing the angular direction of particles following a collision, the coordinates \((\eta, \phi)\) are most often used because the differential production rates of particles are expected to be constant in \( \phi \) and approximately constant in \( \eta \). It is for the same reason that distances between particles, as well as cone sizes used when reconstructing objects from calorimetric clusters or when counting activity near a reconstructed object, are most commonly expressed in terms of \( \Delta R = \sqrt{\Delta \eta^2 + \Delta \phi^2} \).

![Figure 4.8: The ATLAS detector [12].](image-url)
4.3 Inner Detector

The ATLAS Inner Detector, illustrated in Figures 4.10 and 4.11, is designed to identify tracks from charged particles and to reconstruct primary and secondary vertices using these tracks. Closest to the beam pipe, it comprises the Pixel detector, the Semiconductor Tracker (SCT) made of silicon micro-strips, and the Transition Radiation Tracker (TRT), which makes use of straw proportional drift tubes. The Pixel and SCT cover the range $|\eta| < 2.5$, while the TRT covers $|\eta| < 2.0$.

The tracks from charged particles are bent using a solenoid magnetic field. As a result, the radius of curvature of the tracks is proportional to the transverse momentum of the corresponding particles, and their charge is identified using the direction of curvature. The resolution of the ID for measuring the momentum of tracks is discussed in detail in Section 5.2.

Figure 4.9: Illustration of the ways in which different particles are identified based on their interactions with the ATLAS detector [33].
Figure 4.10: Transverse view of the ATLAS Inner Detector components [12].

Figure 4.11: Longitudinal view of the ATLAS Inner Detector components [12].
4.3.1 Pixel Detector

The innermost part of the ATLAS detector is the Pixel detector, which is made of 1744 identical silicon pixel sensor modules, each with 46080 readout channels, for a striking total of more than 80 million readout channels: about 50% of the number for the whole detector. Such a number is necessary to provide the very fine detector granularity allowing to meet the stringent vertex reconstruction requirements imposed by the LHC collision rates.

The dimensions of the pixels are $50 \mu m \times 400 \mu m$ in 90% of cases, and $50 \mu m \times 600 \mu m$ otherwise, with a thickness of 250 $\mu m$. They are arranged as three cylinders in the Barrel region; the first pixel layer, called the $B$-layer for its particular usefulness in reconstructing secondary vertices from $B$-meson decays, is closest to the beryllium beam pipe at a radius of 50.5 mm. On each layer, the modules are laid out in an overlapping way for redundancy, with the top of each module facing the beam pipe in order to minimize the material in front of the detector. In each Endcap region, as illustrated in Figure 4.11, the modules are arranged in three wheels perpendicular to the beam axis. There the redundancy is achieved by alternating the sensors on each side of the supporting structure.

When a charged particle passes through a semiconductor detector, it ionizes atoms in the lattice, resulting in the production of electron-hole pairs. These electrons and holes drift in opposite directions, resulting in a measurable current at the electrodes. The pixels are made of radiation-tolerant doped silicon, operating at a reverse bias voltage of 150 V initially. This voltage is foreseen to be increased up to 600 V after 10 years of operation, to compensate for radiation damage.

An additional inner layer for the Pixel detector, called the Insertable $B$-Layer [14], was inserted during the shutdown between Run-I and Run-II, to further enhance the vertex reconstruction capabilities of the ID. This new innermost pixel layer fits between the other pixel layers and a new, smaller beam pipe.

4.3.2 Semiconductor Tracker

The SCT also makes use of semiconductor detector technology, with 15912 sensors of 768 active micro-strips each. The silicon micro-strips measure $80 \mu m \times 12$ cm, with a thickness of 285 $\mu m$. Their general layout is the same as for the Pixel detector, with four cylindrical layers in the Barrel and nine wheels in each Endcap. The layers and wheels are instrumented on both sides, and a stereo angle of 40 mrad between micro-strips on each side allows 3D measurements of charged particle tracks.

For best operating conditions, both silicon detectors are operated at a temperature between $-5$ and $-10$ degrees Celsius. A laser interferometric monitoring system, able to monitor shape deformations due to temperature variations to an accuracy of 10 $\mu m$, is used to maintain the alignment of the SCT. Like the Pixel detector, the SCT is operated at a voltage of 150 V initially, foreseen to be increased up to 350 V after 10 years of operation to maintain a good charge collection efficiency.
4.3.3 Transition Radiation Tracker

The outermost part of the ID is the TRT, which makes use of straw proportional drift tubes to detect charged particles. The straws have a diameter of 4 mm, a compromise between speed of response and number of hits per track on one hand, and hit efficiency on the other. Indeed while smaller straws are preferable for the first two requirements, with too small a diameter not enough ionization would take place in any straw. It is also important for tracking detectors to have high mechanical stability and rigidity while minimizing the amount of material used. As illustrated in Figure 4.10, the straws in the Barrel are arranged in triangular clusters parallel to the beam axis. In the Endcaps, on the other hand, they form wheels perpendicular to the beam axis, as shown in Figure 4.11.

When charged particles pass through a straw, they ionize the gas it contains; the ions produced drift towards the cathode straw wall, which is kept at a negative tension of $-1530 \text{ V}$, while the corresponding electrons drift towards a gold-plated tungsten anode wire with a small diameter of 31.5 $\mu\text{m}$. When electrons come close to the anode wire, an avalanche occurs with a gain of $2.5 \times 10^4$, resulting in a signal proportional to the energy loss of the ionizing particle. The maximum drift time, corresponding to ionization on the edge of a straw, is measured to be 50 ns, which is satisfactory for operation at LHC collision rates. The drift-time accuracy of the straws is 130 $\mu\text{m}$.

One distinguishing feature of the TRT is its use of transition radiation to distinguish electrons from more massive particles. Transition radiation is emitted when a charged particle crosses the boundary between two materials with different dielectric constants. This radiation is most abundant for highly-relativistic particles with Lorentz factor $\gamma > 10^3$, and in this case the energy of the emitted photons ranges from a few to tens of keV. To exploit this effect, the straw wall consists mainly of two 25-$\mu\text{m}$ kapton films, separated by a 5-$\mu\text{m}$ polyurethane layer: these materials are chosen for their low density and relatively low atomic number $Z$. Indeed, since the photoelectric absorption cross section per atom goes roughly as $Z^5$ [52], a transition radiator with low $Z$ keeps self-absorption small and thus maximizes yield. On the other hand, it important for the gas mixture to absorb the transition radiation photons: xenon is used for this purpose. The straw tubes are filled with a Xe–CO$_2$–O$_2$ mixture at $\{70\%, 27\%, 3\%\}$: CO$_2$ ensures high drift velocities and photon-quenching, and O$_3$, which is created from O$_2$ during avalanches, prevents silicon and hydrocarbon deposits on the anode wire. Unlike the silicon detectors and the liquid argon calorimeters, the TRT operates at room temperature, which requires the presence of heaters.

Two signal thresholds are used by the readout system: the low threshold is triggered by the signal due to ionization energy loss from charged particles, and the high threshold, used for electron identification, is triggered by the significantly stronger signal due to transition radiation. Figure 4.12 shows the average probability of a high-threshold hit in the TRT Barrel as a function of the Lorentz $\gamma$-factor, as measured in the ATLAS combined test-beam. Electrons have a much higher high-threshold probability than pions, while muons are on the rise of the curve. Transition radiation thus allows the TRT to successfully distinguish electrons from pions.
4.3.4 Solenoid Magnet

The superconducting solenoid magnet of ATLAS provides a magnetic field of 2 T parallel to the beam axis. It consists of a single-layer coil made of aluminium-stabilized niobium-titanium, operated at a nominal current of 7.73 kA and a temperature of 4.5 K. This choice of material and layout minimizes the amount of material in front of the calorimeters. The magnetic field flux is returned by the steel of the Tile Hadronic Calorimeter.

4.4 Calorimeters

The ATLAS calorimeters are illustrated in Figure 4.13. Three cryostats, for the Barrel and the two Endcaps, house the calorimeters using liquid argon as the active material: the three parts of the Electromagnetic Calorimeter, the Hadronic Endcap Calorimeters and the Forward Calorimeters. They are operated at a temperature of 88 K, ensuring that the argon remains liquid. The cryostats are surrounded by the Tile Hadronic Calorimeter, which extends above both the Barrel and Endcaps.

4.4.1 Electromagnetic Calorimeter

The Electromagnetic Calorimeter makes use of lead absorbers between layers of liquid argon with kapton electrodes. An accordion shape is used, ensuring a complete symmetry in the $\phi$ coordinate without cracks in coverage. When charged particles traverse the lead absorbers, they emit photons via bremsstrahlung; these photons in turn convert into electron-positron pairs, which emit more photons, the result being a shower of electrons and photons which continues until all the energy is

\[ \text{Lorentz gamma factor} \]

\[ 10^2, 10^3, 10^4, 10^5 \]

\[ \text{High-threshold probability} \]

\[ 0, 0.02, 0.04, 0.06, 0.08, 0.1, 0.12, 0.14, 0.16, 0.18, 0.2, 0.22 \]

Pions, Muons, Electrons

\[ \text{Pions}, \text{Muons}, \text{Electrons} \]

\[ \text{ATLAS} \]

Figure 4.12: Average probability of a high-threshold hit in the TRT Barrel as a function of the Lorentz $\gamma$-factor for electrons (open squares), muons (full triangles) and pions (open circles) in the energy range from 2 to 350 GeV [12].
absorbed by the calorimeter. Photons may also initiate such a shower, with an initial conversion. Particles from the shower ionize the liquid argon, and the ions then drift toward the electrodes resulting in a signal proportional to the energy deposited.

The thickness of the lead absorbers and the granularity of the instrumentation was optimized to provide the best energy and position resolution possible for electrons and photons, especially in the precision region of $|\eta| < 2.5$ corresponding to the acceptance of the ID. The design energy resolution is $\sigma_E/E = 10%/\sqrt{E} \oplus 0.7%$. The coverage of this detector subsystem otherwise extends up to $|\eta| = 3.2$. Showers from photons and electrons are contained in the Electromagnetic Calorimeter, while showers from heavier particles reach the Hadronic Calorimeters.

Figure 4.13: The ATLAS calorimeters [12].

4.4.2 Hadronic Calorimeters

The Hadronic Endcap Calorimeters, located behind the corresponding parts of the Electromagnetic Calorimeter, are comprised of 32 identical wedge-shaped modules each. They make use of copper plates as absorbers, interleaved with gaps of liquid argon and instrumentation, and cover the range $1.5 < |\eta| < 3.2$.

Between the Endcap Calorimeters and the beam pipe, Forward Calorimeters cover the range $3.1 < |\eta| < 4.9$. Instrumenting this region is especially challenging due to the high particle luminosity coming from the collisions. Three modules are present in each Endcap: the first is made of copper and is optimized for electromagnetic measurements, and the others are made of tungsten to measure the intense hadronic interactions in the most compact way possible. Each of these modules
is designed as a metal matrix, with hollow tubes instrumented with metal rods parallel to the beam axis. The very small gaps between the rods and the tube are filled with liquid argon.

Around the three cryostats, the Tile Hadronic Calorimeter covers the range $|\eta| < 1.7$. Its name comes from the scintillating tiles used as the active material, while steel is used as the absorber. Scintillators absorb the energy of charged particles and re-emit it in the form of photons; these photons then travel via fibre optics to readout photomultiplier tubes located around the calorimeter.

The design energy resolution of the Hadronic Calorimeters is $\sigma_E/E = 50%/\sqrt{E} \pm 3\%$, except for the Forward Calorimeter which is limited by its high-radiation operating environment to $\sigma_E/E = 100%/\sqrt{E} \pm 10\%$. Overall, the Hadronic Calorimeters are as hermetic as possible, covering the entire range $|\eta| < 4.9$ and limiting the amount of radiation reaching the Muon Spectrometer.

### 4.5 Muon Spectrometer

The Muon Spectrometer is the most important detector subsystem for the analysis presented in this dissertation, as it is used for the identification and precise momentum measurement of muon tracks. After passing through the calorimeters, muon tracks are bent along the beam direction using a toroidal magnetic field. Muon Drift Tubes (MDTs) measure the track curvature in this direction, in both the detector Barrel and Endcap regions. Cathode Strip Chambers (CSCs) are also present in the very forward region of the detector to extend the muon tracking capability. The muon trigger system comprises Resistive Plate Chambers (RPCs) and Thin Gap Chambers (TGCs); these chambers also provide measurements of the $\phi$ track coordinate. The resolution of the MS for measuring the momentum of charged particle tracks, especially in the case of muons at very high momentum, is discussed in detail in Sections 5.2 and 5.3.2.

#### 4.5.1 Toroid Magnets

The eight-fold symmetric toroid magnets of ATLAS are its most distinctive feature. The magnet coils are made of aluminium-stabilized niobium-titanium-copper, operated at a nominal current of 20.5 kA and a temperature of 4.6 K, producing magnetic fields of approximately 0.5 T in the Barrel and 1 T in the Endcaps. In the Barrel, the coils are encased in stainless-steel vacuum vessels, while in the Endcaps the structure is made of aluminium.

Figure 4.15 shows the integrated magnetic field strength in the MS as a function of $|\eta|$, for $\phi = 0$ and $\phi = \pi/8$. The variability of these field integral values, especially around the magnetic transition region between the Barrel and Endcap toroids, illustrates the importance of using a dependable magnetic field map for tracking muons. It also demonstrates that reliable track momentum measurements require adequate measurements of both angular track coordinates $\eta$ and $\phi$. Momentum measurements are particularly challenging in the transition region, where the magnetic field integral becomes small and even briefly changes sign.
Figure 4.14: The ATLAS Muon Spectrometer and toroid magnets [12].

Figure 4.15: Integrated magnetic field strength in the MS as a function of $|\eta|$, for $\phi = 0$ and $\phi = \pi/8$ [13].
4.5.2 Monitored Drift Tubes

The MDT chambers constitute most of the precision tracking detectors in the MS. The drift tubes are made of aluminium with a diameter of 29.97 mm, and are filled with an Ar–CO$_2$–H$_2$O mixture at {93%, 7%, 300 ppm}, at a pressure of 300 kPa. Muons passing through a tube ionize the gas, and the resulting electrons drift towards a central tungsten-rhenium anode wire with a diameter of 50 $\mu$m, kept at a positive tension of 3080 V. When the electrons come close to the wire, an avalanche occurs with a gain of $2 \times 10^4$, resulting in a measurable current. Because of the larger size of the tubes and the smaller CO$_2$ gas concentration compared to the TRT, the maximum drift time in the MDTs is about 700 ns. Nevertheless, this is satisfactory because the track rates are significantly lower in the MS than in the ID, due to the shielding provided by the calorimeters: at the design LHC luminosity, the maximum counting rate measured by MDTs is expected to be 30 kHz.

Figure 4.16 shows a schematic diagram of the cross section of an MDT, as well as the mechanical structure of an MDT chamber. The chambers are rectangular in the Barrel, and trapezoidal in the Endcaps. They are made of an aluminium frame carrying two multi-layers of three or four drift tube layers each. In each chamber, an internal optical alignment system performs to a precision of 10 $\mu$m, monitoring potential deformations due to temperature changes. Temperature sensors are present as well to quantify these effects. Barrel chambers placed at various horizontal angles are also subject to gravitational sag, which is corrected using a mechanical system informed by the optical alignment system. Finally each chamber is equipped with two to four magnetic field sensors, to ensure a good mapping of the field. The drift-time resolution of individual MDTs is 80 $\mu$m, translating to an optimal tracking resolution of 30 $\mu$m for MDT chambers. The layout of MDT chambers and the global alignment system of the MS are discussed in Section 4.5.6.

![Figure 4.16](image-url)
4.5.3 Cathode Strip Chambers

The CSCs are multi-wire proportional chambers designed to extend the MS coverage in areas where the counting rates are expected to be too high for MDTs to be successful. Each chamber consists of parallel plates instrumented with metallic strips in orthogonal directions. The 5.0-mm gap between the plates is filled with an Ar–CO$_2$ gas mixture at $\{80\%, 20\%\}$. Anode wires having a diameter of 30 $\mu$m, made of gold-plated tungsten with 3% rhenium, are placed mid-way between the plates at intervals of 2.5 mm and operated at a voltage of 1900 V. Electron avalanches at the wires following the passage of charged particles induce signals in the strips on both sides of the CSC.

The precision coordinate $\eta$ is measured using the strips perpendicular to the wires, with a hit resolution of 60 $\mu$m, and a coarser $\phi$ measurement is provided by the wider strips parallel to the wires, with a hit resolution of about 5 mm. The CSC wire signals are not read out.

4.5.4 Resistive Plate Chambers

The RPCs have two objectives: first and foremost, to provide triggering capability in the Barrel of the MS, and second to complement the MDT chambers by measuring the non-precision coordinate $\phi$ of muon tracks. Trigger chambers must be as fast as possible, yet be able to recognize the multiplicity of tracks and perform a coarse measurement of their momentum.

An RPC consists of two units of two rectangular gaseous parallel electrode-plate detectors, as illustrated in Figure 4.17. In each of these detectors, two resistive plates made of plastic laminate are separated by a 2-mm gap created with insulating spacers placed every 100 mm. The gap is filled with C$_2$H$_2$F$_4$–Isobutane–SF$_6$ at $\{94.7\%, 5\%, 0.3\%\}$. The chambers are operated at a voltage of 9.8 kV, creating an electric field of 4.9 kV/mm: with such a high field, ionizing charged particles passing through the gap instantaneously cause an avalanche of electrons which drift towards the anode plate in only a few nanoseconds. By capacitive coupling, this causes a signal on both sides of the gap, and this signal is read out by metallic strips placed in 23–35 mm intervals on the other side of the resistive plates. The strips are oriented to read out the $\eta$ coordinate on one side and the $\phi$ coordinate on the other.

In this way, a charged particle going through an RPC is measured twice in both $\eta$ and $\phi$, unless it crosses in the 65-mm overlap region between units, in which case four measurements are available for each coordinate. The rate of accidental triggers is reduced by requiring independent coincidences in both the $\eta$ and $\phi$ coordinates, within and between RPCs.
4.5.5 Thin Gap Chambers

A trigger detector with higher granularity is needed in the Endcaps, because of the higher radiation levels and the presence of inhomogeneities in the magnetic field in the transition region. TGCs were adopted as the solution in this case: they are multi-wire proportional chambers with a 2.8-mm gap filled with a gas mixture of CO$_2$–n-pentane at \{55\%, 45\\%\}, providing a gain of $3 \times 10^5$. Gold-plated tungsten wires are placed every 1.8 mm, and are operated at a voltage of 2900 V.

The wires are read out to obtain a measurement of $\eta$, while copper strips perpendicular to the wires provide measurements of $\phi$. As a result of the use of a highly-quenching gas combined with a small wire separation, the time resolution of a TGC is smaller than 25 ns for 99\% of tracks, enabling these detectors to act as trigger chambers.

The TGCs are arranged in units of two or three chambers, and make use of coincidence algorithms like the RPCs. Because n-pentane is highly flammable, each TGC unit is surrounded by a gas envelope with circulating CO$_2$, which is monitored. The gas supplies and voltage are automatically switched off if traces of n-pentane are found in the envelope, indicating the presence of a leak.

4.5.6 Layout of the Muon Spectrometer

The layout in the transverse plane of the MS Barrel chambers is illustrated in Figure 4.18. The chambers are placed in three concentric cylindrical layers, called “stations”, each with eight large and eight small chambers matching the arrangement of the toroid magnets. The alternating large
and small sets of chambers are numbered as “sectors”, starting from sector 1 centred on the $x$-axis at $\phi = 0$, with sector numbers increasing with $\phi$. Thus sector 5 corresponds to the top of the detector, and sector 13 to the bottom. Air is utilized as the medium between the stations, for economical reasons but also to minimize the amount of material traversed by the muons in the MS, and thereby reduce multiple-scattering effects.

The MDT standard chambers are named based on their location and features: the detector region (Barrel or Endcap) is indicated first, followed by the station (Inner, Middle or Outer) and finally the size (Large or Small). These names are often abbreviated as the corresponding three-letter acronyms (two-letter acronyms when the size is not specified). Exceptions to the above are necessary to accommodate the supporting feet of the detector:

- In sectors 11 and 15, the BIL chambers are replaced with two chambers named BIR and BIM;
- In sectors 12 and 14, the BMS chambers are replaced with chambers named BMF, and the BOS chambers are replaced with two chambers named BOF and BOG.

These special MDT chambers have proportions designed to match the geometry of the additional supporting material in this region.

Because the toroid magnetic field bends the trajectories of muons in the longitudinal direction, it is crucial for the MS to precisely measure the $\eta$ coordinate of hits: therefore, the orientation of all MDT chambers is such that the tubes are oriented in the $\phi$ direction. The Middle and Outer MDT chambers cover the range $|\eta| < 2.7$; on the other hand, the Inner MDT chambers cover $|\eta| < 2.0$, with the CSC providing coverage for $2.0 < |\eta| < 2.7$.

The RPCs are located on each side of the BM MDT chambers, on the outside of the BOL chambers and on the inside of the BOS chambers. One TGC double-layer is present on the inside of the EI chambers, one triple-layer is on the inside of the EM chambers, and two double-layers are on the outside of the EM chambers. The RPCs and TGCs thus provide triggering capability and $\phi$ track coordinate measurements for $|\eta| < 1.05$ and $1.05 < |\eta| < 2.4$, respectively. Wherever possible, trigger chambers are installed even where there is no space for MDT chambers, in order to maximize the trigger acceptance.

Figure 4.19 shows a longitudinal view of large and small standard MS sectors, namely sectors 3 and 4 respectively: this allows to see the arrangement of both the Barrel and Endcap chambers. In addition to the EI, EM and EO chambers, two extra sets of MDT chambers are present:

- The Endcap Extra (EE) chambers provide additional coverage in the $\eta$ range between the Barrel and the edge of the EO chambers;
- The Barrel Endcap Extra (BEE) chambers, situated around the Endcap toroid magnet cryostats, provide additional coverage in the region where the magnetic field integral changes sign as shown in Figure 4.15.
Within given sectors, the MDT chambers are subdivided into “towers” defined projectively in the $\eta$ coordinate. There are six well-defined towers in the Barrel, while in the Endcaps the arrangement is more complicated, as illustrated in Figure 4.19. Further, unlike the EIL chambers, the EIS chambers can only begin at $|\eta| = 1.25$ because of the Barrel toroid magnets; two additional BIS chambers, named BIS7 and BIS8, are therefore placed in compensation. The BEE and BIS8 chambers are each made of a single multi-layer of drift tubes, instead of the usual double multi-layer.

The holes in coverage between the large MDT chambers near $|\eta| = 0$ are due to services for the ID and the calorimeters. Services for the ID also require small holes between BIL chambers near $|\eta| = 0.64$, while calorimeter services also require a small hole between BIL chambers, but only in sector 13. Larger holes in the coverage of BML and BOL chambers in sector 13 are due to the necessary elevator shafts at the bottom of the detector. Holes between BMS chambers around $|\eta| = 0$, $|\eta| = 0.42$ and $|\eta| = 0.75$ are needed for elements of the Barrel toroid magnet supporting structure. Finally, there are no central EEL chambers in sector 5 to make room for pumps needed by the Endcap toroid magnets [47].
In addition to the internal optical alignment system present in each MDT chamber, a global alignment system is necessary to monitor the position of each tube, in order to achieve a position resolution of 30 $\mu$m. The BI, BM and BO chambers are connected to their neighbours using alignment sensors, and between stations by projective optical lines corresponding to fixed values of $\eta$. Additional lines connect MDT chambers to the Barrel toroid magnets to provide a reference. In the Endcaps, the alignment is performed in two steps: optical lines corresponding to fixed values of $\theta$ link reference high-precision rulers together, and the EI, EE, EM, EO and CSC chambers are optically connected to these rulers and to neighbouring chambers. Unfortunately, the BEE, BIS7 and BIS8 chambers could not be linked to this global alignment system: this fact is taken into account in the selection criteria for muons at very high momentum in Section 5.3.

An important complement to the optical alignment system is the use of track-based alignment algorithms. In particular, straight tracks recorded from collision data taken with the magnetic field turned off provide key information allowing to further improve the alignment of the MS. Such track-based algorithms are described in more detail in Section 5.2.

There is one way in which the MS layout in Run-I differs from the nominal design described above: most of the EE chambers could not be installed in time. Before 2011, they were only installed in sectors 5 and 13 for both Endcaps, sector 11 for $\eta > 0$ and sector 15 for $\eta < 0$, and were not linked to the alignment system. Figure 4.20 shows the number of detector stations traversed by muons in the MS as a function of $\eta$ and $\phi$: the effect of the EE chambers’ absence is seen by comparing the graph for data collected in 2011 to the design layout.

During the winter shutdown between 2011 and 2012, all the EE chambers for $\eta < 0$ were installed, as well as the ones in sector 3 for $\eta > 0$, and the installed EE chambers were fully integrated into the alignment system. The installation of the remaining EE chambers was completed in the spring of 2013, for operation in Run-II.
Figure 4.19: Layout of large (top) and small (bottom) standard sectors of the Muon Spectrometer, specifically sectors 3 and 4 for $\eta \geq 0$. Straight lines are drawn in increments of 0.1 in the $\eta$ coordinate. Figure credit: F. Bauer.
Figure 4.20: Number of detector stations traversed by muons in the MS as a function of $\eta$ and $\phi$ as of March 2011 (top) compared with the design layout (bottom). Figure credit: A. Ouraou.
4.6 Missing Transverse Momentum

The presence of particles which do not interact with any of the detector components can be inferred using the vectorial sum of the transverse momentum measured by all the detector measurements. Among particles from the Standard Model, only neutrino production can cause real missing transverse momentum. On the other hand, many hypotheses beyond the Standard Model predict new particles that cannot be detected directly. Conservation of momentum in the directions transverse to the beam axis implies that the overall transverse momentum vectorial sum from the objects detected after a collision would be zero in the absence of undetectable particles or mis-measurements. The same cannot be said from the longitudinal component of the sum, because the colliding partons have different momenta along the beam axis.

Missing transverse momentum is therefore defined as a 2D vector \([21, 22, 28]\). The magnitude of this vector is commonly called missing transverse energy \((E_{\text{miss}}^T)\), and by an abuse of notation this symbol is also used for its components:

\[
E_{\text{miss}}^i = E_{\text{miss}, e}^i + E_{\text{miss}, \gamma}^i + E_{\text{miss}, \tau}^i + E_{\text{miss}, \text{jets}}^i + E_{\text{miss}, \mu}^i + E_{\text{miss, soft}}^i
\]

with \(i \in \{x, y\}\), where each term of the sum is the negative momentum sum over the calibrated reconstructed objects indicated:

\[
E_{\text{miss, } \alpha}^i = -\sum p_i^\alpha
\]

The order of the terms in Equation 4.1 reflects the priority with which calorimeter energy depositions are associated to objects: electrons first, then photons, hadronic \(\tau\)-lepton decays, jets with \(p_T > 20\) GeV and finally muons. The soft term comprises the calorimeter deposits from jets with \(p_T < 20\) GeV and from unassociated clusters with significant signal. In addition to calorimeter deposits, the momentum of muon candidates entering the MS is also taken into account in the muon term, and the momentum of tracks with low \(p_T\) detected by the ID but missed by the calorimeters is included in the soft term.

4.7 Luminosity Monitoring

A measure of the number \(N_{\text{tot}}\) of collisions having happened inside the detector during a certain period of time is called the integrated luminosity \(L_{\text{int}}\), defined as the integral of the instantaneous luminosity \(L\):

\[
L_{\text{int}} = \int L \, dt = \frac{N_{\text{tot}}}{\sigma_{\text{tot}}}
\]

where \(\sigma_{\text{tot}}\) is the total inelastic proton-proton collision cross section, shown in Figure 2.7. The progression of the integrated luminosity during Run-I of the LHC was shown in Figure 4.5.

Many experimental techniques are used to evaluate this quantity [27]. The main technique used
for calibrating the luminosity measurements outlined below is the van der Meer scan [165], which measures the horizontal and vertical beam widths $\Sigma_x$ and $\Sigma_y$ allowing to calculate $L$ using

$$L = \frac{n_b f_r n_1 n_2}{2\pi \Sigma_x \Sigma_y}$$ (4.4)

where $n_b$ is the number of proton bunches crossing at the interaction point, $f_r$ is the revolution frequency, and $n_1$ and $n_2$ are the numbers of particles in the colliding bunches.

Then, it is possible to estimate the luminosity by counting the number $N_{\text{evt}}$ of collision events detected. These event-counting methods are calibrated during van der Meer scans, and can then be used during regular collision data collection to estimate the integrated luminosity. The detection of tracks in the ID and of energy clusters in the calorimeters are used for this purpose, in addition to the information from dedicated detectors:

- The Beam Condition Monitor, consisting of diamond sensors located near the beam pipe at $|\eta| = 4.2$ and $z = 184$ cm on each side of the interaction point, whose principal function is to protect the detector by providing a fast abort signal in the event of large beam losses;
- The Minimum Bias Trigger Scintillators, located at $2.09 < |\eta| < 3.84$ and used to trigger a sample of events with minimal collision activity;
- The LUCID detector, covering $5.6 < |\eta| < 6.0$, which detects Cherenkov photons produced when charged particles from collisions traverse its aluminium tubes filled with $C_4F_{10}$ gas;
- The Zero Degree Calorimeter, designed to detect very forward neutral particles at $|\eta| > 8.3$ using tungsten absorbers with embedded quartz rods read out by photomultiplier tubes.

As an alternative to van der Meer scans, the ALFA detector, located about 240 m from the ATLAS interaction point, uses scintillators to measure elastic proton-proton scattering rates, which can also in principle calibrate the luminosity measurements. The placement of the LUCID, Zero Degree Calorimeter and ALFA detectors along the LHC beam line around the ATLAS interaction point is shown in Figure 4.4.

From the number of events detected, it is possible to obtain the average number $\lambda$ of simultaneous proton-proton interactions per bunch crossing. Using Poisson statistics, the probability for at least one interaction to occur during a given bunch crossing is

$$\frac{N_{\text{evt}}}{N_{\text{BC}}} = 1 - e^{-\lambda}$$ (4.5)

and therefore

$$\lambda = -\ln \left(1 - \frac{N_{\text{evt}}}{N_{\text{BC}}}\right)$$ (4.6)
where $N_{BC}$ is the total number of bunch crossings having happened during the time $N_{\text{ev}}$ was measured. It follows that the instantaneous luminosity is

$$L = \frac{\lambda m f_{\text{tr}}}{\sigma_{\text{vis}}}$$

where the visible cross section $\sigma_{\text{vis}}$ is the quantity calibrated during van der Meer scans by comparing this result to the value measured using Equation 4.4.

In practice, the formalism is more involved, taking into account bunch-by-bunch differences in instantaneous luminosity and corrections due to pileup in the detector. As well, coincidence algorithms are used in addition to the simple counting algorithms described above in order to evaluate the impact of non-collision backgrounds. The total integrated luminosity was calculated to a precision of 1.8% for data collected in 2011 [27], and the corresponding preliminary uncertainty value for data collected in 2012 is 3.6%.

It is also possible to convert the measured rates for the production of $W$ and $Z$ bosons into a measure of the integrated luminosity using the theoretical cross section for these processes. While this last technique is not used in general, in order to retain the possibility to carry out $W$ and $Z$ boson production cross section measurements at ATLAS, it is effectively used in the search presented in this dissertation, which normalizes the total background prediction to the measured event yields on the $Z$ peak. This is further explained in Sections 6.2.1 and 7.1.

### 4.8 Trigger, Data Acquisition and Data Quality

#### 4.8.1 Trigger

As is readily seen from Figure 2.7 in Section 2.2.2, the different production cross sections at the LHC span many orders of magnitude, with processes involving $b$-quarks ranking factors 100 to 1000 smaller than the total cross section, electroweak boson production another factor $10^4$ smaller and Higgs boson production yet another factor $10^3$ to $10^5$ smaller. Processes beyond the Standard Model are also expected to have very small rates compared to the total proton-proton inelastic cross section. In parallel, the available data collection bandwidth and storage capacity of ATLAS is significantly smaller than the enormous event rate. A reliable and efficient trigger system is therefore crucial to the success of the experiment, in order to select the collision data that contains physical processes of interest among very large backgrounds.

ATLAS makes use of a dedicated trigger system that consists of three layers, respectively called Level-1, Level-2 and the Event Filter. The Level-1 trigger is a hardware-based trigger, which selects collision events of interest based on features readily available during early reconstruction. Specifically, it uses information from the trigger chambers of the MS and reduced-granularity information from all the calorimeters to look for high-$p_T$ muons, electrons, photons, hadronic $\tau$-lepton decays...
and jets, as well as large $E_T^{\text{miss}}$. This first step, with an allocated decision time of $2.5 \, \mu s$, already makes possible a strong reduction of the event rate from $20 \, \text{MHz}$ (for a proton bunch time separation of 50 ns) to $\mathcal{O}(100 \, \text{kHz})$.

For each event selected by the Level-1 trigger, one or more “regions of interest” are defined, with information about the preliminary angular coordinates of interesting objects and the Level-1 threshold passed. This information is passed to the two higher trigger levels, which use software algorithms in order to determine which of these events are to be recorded for offline processing. The Level-2 trigger fully reconstructs all the event data from the region of interest identified at Level-1, which represents approximately $2\%$ of the total event data, with an average event processing time of $40 \, \text{ms}$ to further reduce the event rate to about $5 \, \text{kHz}$. Then the Event Filter reconstructs the complete event in an average time of $4 \, \text{seconds}$, finally reducing the rate to $400 \, \text{Hz}$, including a $5\text{-Hz}$ random trigger rate. Considering an average of about 20 interactions per crossing at $20 \, \text{MHz}$ in 2012, this means that out of all the proton-proton collisions observed by the ATLAS detector, only one out of a million is stored for analysis.

In fact, some interesting physical processes have cross sections that are too large to allow the storage all the produced events. This implies that the corresponding triggers have to be prescaled, meaning that only one in a number of events passing the triggers are kept. For example, in Run-I, events with a $J/\psi$ meson quickly became too numerous to allow keeping them all; for Run-II this situation is even foreseen to affect $W \rightarrow \ell \nu$ production. Nevertheless, in order to maximize the signal sensitivity to new, rare physical processes, searches strive to make use of un-prescaled triggers whenever possible.

## 4.8.2 Data Acquisition

Each detector channel in ATLAS is linked to front-end electronics, which perform signal digitization and provide buffers long enough to make possible the operation of the Level-1 trigger. Following the Level-1 decision, the selected events are stored in other buffers until they are transferred via readout drivers to the Data Acquisition (DAQ) system, where the Level-2 and Event Filter trigger algorithms are run and monitored.

The DAQ system is configured for every data-taking period, or “run”, with a specific trigger menu corresponding to the objectives of the run. Each run is further subdivided into 1- or 2-minute intervals called “luminosity blocks”: this segmentation allows the DAQ system to control the data-taking conditions without having to re-start the run, which would involve data losses.

A hardware monitoring system called the Detector Control System (DCS) ensures that any abnormal behaviour in the detector is reported to operators: automatic or manual corrective actions are then taken. Example quantities monitored by the DCS include temperature, humidity, gas concentration and pressure values, magnetic field, voltage, etc. The DCS also handles the communication between ATLAS subsystems, as well as between ATLAS and external systems at CERN.
4.8.3 Data Quality

Similar to the fact that only a fraction of the collisions delivered to ATLAS are recorded by the detector, only a fraction of the recorded collision events are suitable for use in analyses. As shown in Figure 4.5, this fraction is high: in 2011, 4.57 fb$^{-1}$ out of 5.08 fb$^{-1}$ (90.0%) of the recorded collision data passing all quality requirements; for 2012 this proportion is even higher at 20.3 fb$^{-1}$ out of 21.3 fb$^{-1}$ (95.3%).

The quality of the data is first verified as it is being recorded: this is called “online monitoring”. A subset of the recorded events, called the Express Stream, is drawn and reconstructed in real-time. Many data quality verifications are performed automatically. In addition, data quality shifters in the ATLAS Control Room manually monitor a series of histograms corresponding to a wide variety of global quantities essential for the successful operation of the detector. For example, it is necessary to verify the synchronization of all detector subsystems with each other, as well as distributions of the number of hits observed in each subsystem, detector occupancy maps, magnetic field measurements, correlations between the parameters of the tracks observed in both the ID and the MS, etc.

Physical objects are also monitored: electron, muon, and $\tau$ lepton candidates, as well as hadrons and missing transverse momentum. It is even possible to reconstruct $J/\psi$ and $\Upsilon$ mesons from events with pairs of opposite-sign muon candidates, and $W$ and $Z$ bosons in the electron and muon channels, immediately as these objects are detected by ATLAS. Since the successful observation of these objects relies on the good performance of all detector subsystems, monitoring their observation rates as data are collected is an effective way to become aware of new detector problems as they occur. Indeed, online data quality monitoring shifters are often the first to identify such problems; this information is then transmitted to the shifters dedicated to the relevant subsystems.

My main contributions to the ATLAS data collection were made in the context of this framework, by developing monitoring histograms and automatic algorithms, maintaining the software packages and helping to define the instructions for data quality shifters in the ATLAS Control Room.

Additional data quality verifications, called “offline monitoring”, are performed after the complete reconstruction of the data. They involve essentially the same quantities as in the case of online monitoring, in addition to any validation which requires high statistics from complete runs. When detector problems are found, this information is stored in the data quality database [109], indicating which luminosity blocks may not be used for certain purposes.

The determination of the dataset suitable for analysis is finalized when dedicated lists of good luminosity blocks are prepared from the data quality database, catering to the specific needs of the different analysis categories. For instance, the search for new neutral resonances decaying into muon pairs presented in this dissertation requires good data quality from the ID and from the MS. On the other hand, this analysis is insensitive to a number of data quality problems occurring in the calorimeters, and thus benefits from a slightly higher integrated luminosity compared to analyses for which all data quality requirements are necessary.
Chapter 5

Muons at Very High Momentum in ATLAS

In order to successfully reconstruct an eventual resonant peak at dimuon invariant masses above the Z peak, while preventing a potential contamination of the signal region by false high-momentum tracks originating from mis-reconstruction, a detailed experimental understanding of muons at very high momentum is absolutely necessary. Only the muon candidates traversing the regions of ATLAS offering the best momentum resolution can be selected. Section 5.1 gives an overview of the different types of muon objects reconstructed in ATLAS, and Section 5.2 details the performance of the type of muon candidates used in this search.

The last section explains the exact muon candidate selection used in this analysis. Its first part, Section 5.3.1 concerns muon candidates with segments from all three stations of the Muon Spectrometer (MS), called “3-station muons” for short: these are the best-reconstructed muon tracks in ATLAS. The muon momentum measurement is derived from the curvature of the track, which for 3-station muons is obtained from the sagitta, that is the distance between the middle track segment and a straight line linking the inner and outer segments in the MS. The calculation of the sagitta is illustrated in Figure 5.1.

While for this analysis the primary concern is to ensure that all selected muon tracks are well-measured, it is also very important to maximize the signal acceptance of the search, defined as the number of selected signal events over the total number present in the collision dataset. In order to increase the signal acceptance in this analysis, the best muon candidates with two MS track segments are also considered: these are called “2-station muons”. In this case the track curvature is measured using the angular difference between the two measured segments. Dedicated studies are necessary in order to ensure that the resolution of the selected 2-station muons is sufficient for use in this search: these studies are detailed in Section 5.3.2.
5.1 Muon Reconstruction in ATLAS

Four different types of muon candidates are reconstructed from the raw ATLAS data, according to the information available in the different subsystems [13, 32, 35]. The main information consists of the charged particle track measurements, performed independently in the ID and the MS. In the ID, the measurement comes from a fit to hits detected in the Pixel, SCT and TRT detectors.

In the MS, the hits from the chambers in individual stations are first used to form track segments, which are then combined to form a track. These MS tracks extrapolated to the primary vertex of the event are called standalone muons. Energy depositions in the calorimeters are also taken into account when extrapolating the tracks measured in the MS to the primary vertex of the event, by adding the energy losses of the muon candidate to the measurement from the MS. The methods used to estimate these energy losses are explained in detail in Ref. [139].

Combined muons are the best-measured muon candidates from the detector. They are formed when a standalone muon track is successfully matched to a track in the ID. This is only possible within the acceptance of the ID, that is for detector $|\eta| < 2.5$, while only standalone muons are available in the range $2.5 < |\eta| < 2.7$. During Run-I of the LHC, two independent algorithms have been used to reconstruct combined muons. One of them, named “Staco” or “chain 1”, performs a statistical combination of the standalone and ID muon track parameters using the covariance matrices of the track parameter measurements from each subsystem. The other, named “Muid” or “chain 2”, obtains combined muons by performing a global refit from the hits measured in both the ID and the MS. This analysis makes use of the Muid algorithm.

The last two types of muon objects, segment-tagged muons and calorimeter-tagged muons, are useful in analyses prioritizing muon acceptance. In both cases the expected muon momentum has to be small enough for the measurement from the ID to be reliable, as no momentum measurement is available from the MS. Segment-tagged muons consist of tracks measured in the ID that are

![Figure 5.1: Determination of the sagitta based on the location of segments in the Muon Spectrometer. Figure credit: P.-F. Giraud.](image)
matched to single track segments in the MS. They can be used to recover muons in a few areas between the Barrel and Endcaps of the detector ($1.1 < |\eta| < 1.3$), where only one MS station is installed. Finally, services for the ID and the calorimeters pass in the most central region of the detector for $|\eta| < 0.1$, implying the existence of multiple areas where no MS chamber could be installed. Calorimeter-tagged muons, consisting of ID tracks matched to energy depositions in the calorimeter where these depositions are consistent with the ones expected from a minimum-ionizing particle, can be used to recover acceptance in this region.

5.2 High-Momentum Muon Performance

In this analysis, the muons at very high momentum expected from high-mass resonances require precise momentum measurements from the detector, and while signal acceptance is also very important, it comes as a secondary concern. Only combined muons are therefore selected in this search.

The momentum of muons is measured using the curvature of their tracks caused by the solenoid magnetic field in the ID and the toroid magnetic field in the MS. Since this curvature is inversely proportional to the muon momentum, the tracks of muons at very high momentum are very straight, and the uncertainty on the momentum measurement is dominated by the intrinsic resolution and alignment of the Muon Spectrometer. Specifically, this uncertainty can be parametrized as follows:

$$\frac{\sigma(p_T)}{p_T} = \frac{P_0}{p_T} \oplus P_1 \oplus P_2 \cdot p_T$$

(5.1)

where $P_0$ is the resolution parameter related to energy loss fluctuations, $P_1$ is related to multiple scattering, and $P_2$ is related to the alignment and intrinsic resolution. Figure 5.2 illustrates these different contributions to the muon momentum resolution as a function of muon $p_T$, as documented at the time of the Muon Spectrometer Technical Design Report [11]. The design resolution at very high momentum is $\sigma(p_T)/p_T = 10\%/\text{TeV}$. Given the intrinsic resolution described in Section 4.5, achieving this resolution for 3-station muons requires an alignment with sagitta bias lower than 40 $\mu$m [107].

Later estimates for the resolution parameters, which are used in simulated samples for the analyses at $\sqrt{s} = 7$ TeV and $\sqrt{s} = 8$ TeV, are respectively shown in Tables 5.1 and 5.2. They are determined from the known geometry of the ATLAS detector including the distribution of material, the number of hits on each track and the intrinsic hit resolution, as well as magnetic field integrals and the expected resolution of the detector as determined from the optical alignment system whenever possible. Most resolution parameters are taken to be constant in regions of $\eta$. The exception is the $P_2^{ID}$ term for $2.0 < |\eta| < 2.5$: since the TRT coverage ends at $|\eta| = 2.0$, the $P_2^{ID}$ term, which is dependent on the muon track length in the detector active material, is proportional to $\sinh^2 \eta$ in this region.
Table 5.1: Muon $p_T$ resolution parameters for simulated samples generated in 2011 at $\sqrt{s} = 7$ TeV.

<table>
<thead>
<tr>
<th>Region</th>
<th>$p_1^{ID}$ [%]</th>
<th>$p_2^{ID}$ [TeV$^{-1}$]</th>
<th>$p_1^{MS}$ [%]</th>
<th>$p_2^{MS}$ [TeV$^{-1}$]</th>
</tr>
</thead>
<tbody>
<tr>
<td>$</td>
<td>\eta</td>
<td>&lt; 1.05$</td>
<td>1.61</td>
<td>0.31</td>
</tr>
<tr>
<td>$1.05 &lt;</td>
<td>\eta</td>
<td>&lt; 1.7$</td>
<td>2.59</td>
<td>0.33</td>
</tr>
<tr>
<td>$1.7 &lt;</td>
<td>\eta</td>
<td>&lt; 2.0$</td>
<td>3.39</td>
<td>0.44</td>
</tr>
<tr>
<td>$2.0 &lt;</td>
<td>\eta</td>
<td>&lt; 2.5$</td>
<td>5.12</td>
<td>0.042 sinh$^2 $\eta</td>
</tr>
</tbody>
</table>

Table 5.2: Muon $p_T$ resolution parameters for simulated samples generated in 2012 at $\sqrt{s} = 8$ TeV.

<table>
<thead>
<tr>
<th>Region</th>
<th>$p_1^{ID}$ [%]</th>
<th>$p_2^{ID}$ [TeV$^{-1}$]</th>
<th>$p_1^{MS}$ [%]</th>
<th>$p_2^{MS}$ [TeV$^{-1}$]</th>
</tr>
</thead>
<tbody>
<tr>
<td>$-2.50 &lt; \eta &lt; -2.25$</td>
<td>4.88</td>
<td>0.0523 sinh$^2 $\eta</td>
<td>3.09</td>
<td>0.208</td>
</tr>
<tr>
<td>$-2.25 &lt; \eta &lt; -2.00$</td>
<td>4.88</td>
<td>0.0523 sinh$^2 $\eta</td>
<td>3.09</td>
<td>0.208</td>
</tr>
<tr>
<td>$-2.00 &lt; \eta &lt; -1.70$</td>
<td>3.30</td>
<td>0.447</td>
<td>3.44</td>
<td>0.208</td>
</tr>
<tr>
<td>$-1.70 &lt; \eta &lt; -1.50$</td>
<td>2.50</td>
<td>0.322</td>
<td>4.77</td>
<td>0.287</td>
</tr>
<tr>
<td>$-1.50 &lt; \eta &lt; -1.05$</td>
<td>2.50</td>
<td>0.322</td>
<td>4.77</td>
<td>0.287</td>
</tr>
<tr>
<td>$-1.05 &lt; \eta &lt; -0.80$</td>
<td>1.58</td>
<td>0.299</td>
<td>3.25</td>
<td>0.188</td>
</tr>
<tr>
<td>$-0.80 &lt; \eta &lt; -0.40$</td>
<td>1.58</td>
<td>0.299</td>
<td>3.25</td>
<td>0.188</td>
</tr>
<tr>
<td>$-0.40 &lt; \eta &lt; 0.00$</td>
<td>1.58</td>
<td>0.299</td>
<td>3.25</td>
<td>0.188</td>
</tr>
<tr>
<td>$0.00 &lt; \eta &lt; 0.40$</td>
<td>1.58</td>
<td>0.299</td>
<td>3.25</td>
<td>0.188</td>
</tr>
<tr>
<td>$0.40 &lt; \eta &lt; 0.80$</td>
<td>1.58</td>
<td>0.299</td>
<td>3.25</td>
<td>0.188</td>
</tr>
<tr>
<td>$0.80 &lt; \eta &lt; 1.05$</td>
<td>1.58</td>
<td>0.299</td>
<td>3.25</td>
<td>0.188</td>
</tr>
<tr>
<td>$1.05 &lt; \eta &lt; 1.50$</td>
<td>2.59</td>
<td>0.316</td>
<td>5.25</td>
<td>0.323</td>
</tr>
<tr>
<td>$1.50 &lt; \eta &lt; 1.70$</td>
<td>2.59</td>
<td>0.316</td>
<td>5.25</td>
<td>0.323</td>
</tr>
<tr>
<td>$1.70 &lt; \eta &lt; 2.00$</td>
<td>3.28</td>
<td>0.450</td>
<td>3.60</td>
<td>0.185</td>
</tr>
<tr>
<td>$2.00 &lt; \eta &lt; 2.25$</td>
<td>4.85</td>
<td>0.0516 sinh$^2 $\eta</td>
<td>3.02</td>
<td>0.219</td>
</tr>
<tr>
<td>$2.25 &lt; \eta &lt; 2.50$</td>
<td>4.85</td>
<td>0.0516 sinh$^2 $\eta</td>
<td>3.02</td>
<td>0.219</td>
</tr>
</tbody>
</table>
Measurements of the resolution of the detector are performed in situ using fits to the invariant mass peaks from $J/\psi \rightarrow \mu^+\mu^-$ and $Z \rightarrow \mu^+\mu^-$ candidate events. Such measurements mainly constrain the $P_0$ and $P_1$ terms from Equation 5.1. Additional constraints on $P_2$ are obtained using straight tracks from collision data recorded in runs with the magnetic field turned off. In particular, sagitta measurements from these straight tracks indicate the sagitta bias of each MS tower. The measured sagitta bias values are found to be within the design resolution for many towers, and within 100 $\mu$m for almost all towers in the Barrel and Endcap, except in the CSC region where biases of up to 250 $\mu$m are observed \[107\].

Muon tracks traversing overlapping sets of chambers in the MS are also used to quantify residual mis-alignments, by comparing the track segments measured by individual chambers in each station. The independent momentum measurements from the overlapping towers can also be compared directly, by fitting the muon track once with the hits in the small sector only, and once with the hits in the large sector only. The quantity \( (q/p_T)_{\text{small}} - (q/p_T)_{\text{large}} \) is then a measure of the momentum bias.

The event with the highest dimuon invariant mass recorded in 2011, displayed in Figure 5.3, includes an excellent example of a muon candidate passing through overlapping MS chambers. It is also worth noticing that the other muon candidate in the event passes the 2-station muon selection. The resolution performance of 2-station muons at very high momentum is described in more detail in Section 5.3.2.

---

\[1\] Recent updates to the MS alignment realized after the analyses presented here were completed have improved the resolution in all towers including in the CSC region, where the sagitta bias values are now within 80 $\mu$m.
Figure 5.3: Event with the highest dimuon invariant mass observed in data collected at $\sqrt{s} = 7$ TeV by the ATLAS experiment. The two muon candidates have transverse momenta of 648 GeV and 583 GeV respectively, and the dimuon invariant mass is 1.25 TeV. The muon candidate in the upper half of the detector is a 2-station muon, and the other muon track passes through overlapping chambers in the MS.
To address differences between the resolution parameters observed in data with respect to the values used in simulations, corrections are applied to the transverse momentum $p_T$ of each simulated muon. This correction, applied independently in the ID and the MS, takes the form of a Gaussian smearing of the quantity $q/p_T$:

$$\delta (q/p_T) = S_1 \cdot g_1 \cdot (q/p_T) + S_2 \cdot g_2$$  \hspace{1cm} (5.2)

where $g_1$ and $g_2$ are random Gaussian variables with zero mean and unit standard deviation, and $S_1$ and $S_2$ are smearing constants. These smearing constants are determined as a function of pseudorapidity by taking the differences in quadrature between the resolution parameters from data and the Monte Carlo simulations:

$$S_i = \begin{cases} P_{i,\text{data}} \ominus P_{i,\text{MC}} & \text{if } P_{i,\text{data}} > P_{i,\text{MC}} \\ 0 & \text{otherwise} \end{cases}$$  \hspace{1cm} (5.3)

The correction to the combined $q/p_T$ is then the weighted average of the individual corrections in the ID and MS:

$$\delta (q/p_T)_{\text{CB}} = \frac{1}{\sigma_{\text{ID}}^2 + \sigma_{\text{MS}}^2} \left( \frac{\delta (q/p_T)_{\text{ID}}}{\sigma_{\text{ID}}^2} \cdot \frac{P_{\text{ID}}}{P_{\text{CB}}} + \frac{\delta (q/p_T)_{\text{MS}}}{\sigma_{\text{MS}}^2} \cdot \frac{P_{\text{MS}}}{P_{\text{CB}}} \right)$$  \hspace{1cm} (5.4)

where $P_{\text{ID}}$, $P_{\text{MS}}$, $\sigma_{\text{ID}}$, and $\sigma_{\text{MS}}$ are the un-smeared transverse momentum values, and $\sigma_{\text{det}} = \sigma(p_{\text{det}})$ are the expected resolutions of the two subsystems calculated using Equation\[5.1\].

The smearing constant values used are shown in Table 5.3 for analyses at $\sqrt{s} = 7$ TeV, and in Table 5.4 for analyses at $\sqrt{s} = 8$ TeV. For the latter, the nominal values of the smearing constants $S_1^{\text{ID}}$ and $S_2^{\text{MS}}$ are set to zero because the resolution parameters used in data turned out to be slightly better than the ones used in the simulation. Upper bounds on these parameters are still considered when propagating systematic uncertainties to the analysis results [32].

**Table 5.3:** Muon momentum smearing constants used in analyses at $\sqrt{s} = 7$ TeV. The effect of $S_1^{\text{ID}}$ is neglected.

<table>
<thead>
<tr>
<th>Region</th>
<th>$S_2^{\text{ID}}$ [TeV$^{-1}$]</th>
<th>$S_1^{\text{MS}}$ [%]</th>
<th>$S_2^{\text{MS}}$ [TeV$^{-1}$]</th>
</tr>
</thead>
<tbody>
<tr>
<td>$</td>
<td>\eta</td>
<td>&lt; 1.05$</td>
<td>$0.19 \pm 0.01$</td>
</tr>
<tr>
<td>$1.05 &lt;</td>
<td>\eta</td>
<td>&lt; 1.7$</td>
<td>$0.24 \pm 0.03$</td>
</tr>
<tr>
<td>$1.7 &lt;</td>
<td>\eta</td>
<td>&lt; 2.0$</td>
<td>$0.50 \pm 0.02$</td>
</tr>
<tr>
<td>$2.0 &lt;</td>
<td>\eta</td>
<td>&lt; 2.5$</td>
<td>$(0.015 \pm 0.004) \sinh^2 \eta$</td>
</tr>
</tbody>
</table>
Table 5.4: Muon momentum smearing constants used in analyses at $\sqrt{s} = 8$ TeV.

<table>
<thead>
<tr>
<th>Region</th>
<th>$S_{1}^{ID}$ [%]</th>
<th>$S_{2}^{ID}$ [TeV$^{-1}$]</th>
<th>$S_{1}^{MS}$ [%]</th>
<th>$S_{2}^{MS}$ [TeV$^{-1}$]</th>
</tr>
</thead>
<tbody>
<tr>
<td>$-2.50 &lt; \eta &lt; -2.25$</td>
<td>0.00 $^{+0.49}_{-0.00}$</td>
<td>$(0.0073 \pm 0.0086) \sinh^{2} \eta$</td>
<td>1.59 $\pm 0.17$</td>
<td>0.00 $^{+0.16}_{-0.00}$</td>
</tr>
<tr>
<td>$-2.25 &lt; \eta &lt; -2.00$</td>
<td>0.00 $^{+0.49}_{-0.00}$</td>
<td>$(0.0289 \pm 0.0041) \sinh^{2} \eta$</td>
<td>1.57 $\pm 0.13$</td>
<td>0.00 $^{+0.16}_{-0.00}$</td>
</tr>
<tr>
<td>$-2.00 &lt; \eta &lt; -1.70$</td>
<td>0.00 $^{+0.17}_{-0.00}$</td>
<td>0.340 $\pm 0.028$</td>
<td>1.64 $\pm 0.15$</td>
<td>0.00 $^{+0.16}_{-0.00}$</td>
</tr>
<tr>
<td>$-1.70 &lt; \eta &lt; -1.50$</td>
<td>0.00 $^{+0.17}_{-0.00}$</td>
<td>0.310 $\pm 0.021$</td>
<td>2.07 $\pm 0.08$</td>
<td>0.00 $^{+0.22}_{-0.00}$</td>
</tr>
<tr>
<td>$-1.50 &lt; \eta &lt; -1.05$</td>
<td>0.00 $^{+0.13}_{-0.00}$</td>
<td>0.275 $\pm 0.016$</td>
<td>1.35 $\pm 0.10$</td>
<td>0.00 $^{+0.22}_{-0.00}$</td>
</tr>
<tr>
<td>$-1.05 &lt; \eta &lt; -0.80$</td>
<td>0.00 $^{+0.13}_{-0.00}$</td>
<td>0.248 $\pm 0.041$</td>
<td>0.34 $\pm 0.26$</td>
<td>0.00 $^{+0.14}_{-0.00}$</td>
</tr>
<tr>
<td>$-0.80 &lt; \eta &lt; -0.40$</td>
<td>0.00 $^{+0.08}_{-0.00}$</td>
<td>0.206 $\pm 0.019$</td>
<td>0.30 $\pm 0.11$</td>
<td>0.00 $^{+0.14}_{-0.00}$</td>
</tr>
<tr>
<td>$-0.40 &lt; \eta &lt; 0.00$</td>
<td>0.00 $^{+0.08}_{-0.00}$</td>
<td>0.229 $\pm 0.013$</td>
<td>0.98 $\pm 0.11$</td>
<td>0.00 $^{+0.14}_{-0.00}$</td>
</tr>
<tr>
<td>$0.00 &lt; \eta &lt; 0.40$</td>
<td>0.00 $^{+0.08}_{-0.00}$</td>
<td>0.208 $\pm 0.016$</td>
<td>1.03 $\pm 0.10$</td>
<td>0.00 $^{+0.14}_{-0.00}$</td>
</tr>
<tr>
<td>$0.40 &lt; \eta &lt; 0.80$</td>
<td>0.00 $^{+0.08}_{-0.00}$</td>
<td>0.203 $\pm 0.016$</td>
<td>0.11 $^{+0.48}_{-0.11}$</td>
<td>0.00 $^{+0.14}_{-0.00}$</td>
</tr>
<tr>
<td>$0.80 &lt; \eta &lt; 1.05$</td>
<td>0.00 $^{+0.13}_{-0.00}$</td>
<td>0.237 $\pm 0.007$</td>
<td>0.53 $\pm 0.11$</td>
<td>0.00 $^{+0.14}_{-0.00}$</td>
</tr>
<tr>
<td>$1.05 &lt; \eta &lt; 1.50$</td>
<td>0.00 $^{+0.13}_{-0.00}$</td>
<td>0.269 $\pm 0.014$</td>
<td>0.83 $\pm 0.15$</td>
<td>0.00 $^{+0.24}_{-0.00}$</td>
</tr>
<tr>
<td>$1.50 &lt; \eta &lt; 1.70$</td>
<td>0.00 $^{+0.17}_{-0.00}$</td>
<td>0.284 $\pm 0.023$</td>
<td>2.15 $\pm 0.16$</td>
<td>0.00 $^{+0.24}_{-0.00}$</td>
</tr>
<tr>
<td>$1.70 &lt; \eta &lt; 2.00$</td>
<td>0.00 $^{+0.17}_{-0.00}$</td>
<td>0.378 $\pm 0.010$</td>
<td>1.53 $\pm 0.09$</td>
<td>0.00 $^{+0.14}_{-0.00}$</td>
</tr>
<tr>
<td>$2.00 &lt; \eta &lt; 2.25$</td>
<td>0.00 $^{+0.49}_{-0.00}$</td>
<td>$(0.0310 \pm 0.0057) \sinh^{2} \eta$</td>
<td>1.46 $\pm 0.14$</td>
<td>0.00 $^{+0.17}_{-0.00}$</td>
</tr>
<tr>
<td>$2.25 &lt; \eta &lt; 2.50$</td>
<td>0.00 $^{+0.49}_{-0.00}$</td>
<td>$(0.0050 \pm 0.0028) \sinh^{2} \eta$</td>
<td>1.22 $\pm 0.21$</td>
<td>0.00 $^{+0.17}_{-0.00}$</td>
</tr>
</tbody>
</table>

The muon momentum scale is also constrained by the fit to the $Z \to \mu^{+}\mu^{-}$ mass peak. As shown in Figure 5.4 for data collected at $\sqrt{s} = 8$ TeV, these corrections are under 0.3% for all simulated muons [32]. Corrections applied in for data at $\sqrt{s} = 7$ TeV are of the same order.

The effect of the corrections to the muon momentum scale and resolution on the $Z$ peak are shown in Figure 5.5. Figure 5.6 shows the impact of the corrections derived from data at $\sqrt{s} = 7$ TeV on a simulated $Z'_{SSM}$ resonant peak with a pole mass of 2.25 TeV, and on the steeply falling background estimate in this region.

The efficiency of the ATLAS detector for reconstructing muons is measured using a sample of $Z \to \mu^{+}\mu^{-}$ events satisfying the requirement $|m_{\mu^{+}\mu^{-}} - M_{Z}| < 10$ GeV, with both muons having $p_{T} > 20$ GeV. The technique used is called “tag-and-probe”: one of the muon candidates, called the “tag”, is required to be a combined muon, while the requirements on the other muon track, called the “probe”, are looser. When measuring the reconstruction efficiency in the ID, the probe muon
is required to be a standalone muon, while ID tracks or calorimeter-tagged muons can be used as probes when measuring the efficiency in the MS.

The reconstruction efficiency in each detector subsystem then corresponds to the proportion of probes matching a second combined muon in the event. Finally, the efficiency for reconstructing combined muons is the product of the track reconstruction efficiency in the ID with the MS reconstruction efficiency given the presence of an ID track. Figure 5.7 shows the combined muon efficiency as a function of pseudorapidity, for data collected at $\sqrt{s} = 8$ TeV and a corresponding simulated $Z \rightarrow \mu^+\mu^-$ sample. In order to correct the muon efficiency in simulated samples to the values observed in data, event-by-event scale factors are applied for each simulated muon.

Trigger efficiency is also estimated using a tag-and-probe technique, but here the measured quantity is the proportion of probes matched to a triggered muon.

**Figure 5.4:** Momentum scale corrections to ID (top) and MS (bottom) tracks derived from data collected at $\sqrt{s} = 8$ TeV [32].
Figure 5.5: Effect of corrections to the muon momentum scale and resolution in the $Z$ peak region. Data and simulated background are shown at $\sqrt{s} = 8$ TeV. The histogram on the top left compares the uncorrected simulation to the data. On the top right, corrections to the momentum resolution are applied, and corrections to both the scale and resolution are applied on the bottom histogram. Additional shape differences are due to the emission of initial-state and final-state radiation which was not fully taken into account in the simulated events used for these histograms [32].
Figure 5.6: Effect of muon resolution smearing on a $Z'_{SSM}$ resonant peak with a pole mass of 2.25 TeV, and on the steeply falling background from $Z/\gamma^*$ production.

Figure 5.7: Efficiency of the ATLAS detector for reconstructing combined muons as a function of pseudorapidity, for data collected at $\sqrt{s} = 8$ TeV and a corresponding simulated $Z \rightarrow \mu^+\mu^-$ sample [32].
5.3 Dedicated Very-High Momentum Muon Selection

5.3.1 Selection of 3-Station Muons

To guard against false high-momentum tracks that could arise from mis-reconstruction, all selected muon candidates must pass stringent quality requirements. Combined muons with segments in all three stations of the Muon Spectrometer constitute the major part of the muons selected in this analysis. The requirements on these 3-station muons for the analysis at \( \sqrt{s} = 8 \) TeV are as follows:

- The measured muon transverse momentum must be above 25 GeV.

- To guarantee the quality of the track in the Inner Detector, each muon must pass the following requirements on the number of detector hits:
  
  - At least one hit in the first layer of the Pixel detector, if such a hit is expected;
  
  - At least 1 hit in the Pixel detector, including Pixel dead sensors crossed;
  
  - At least 5 hits in the SCT, including SCT dead sensors crossed;
  
  - At most 2 Pixel or SCT holes (silicon detector holes are instances where an operational sensor registers no hit but one is expected from the reconstructed track);
  
  - If the muon is within the range \( 0.1 < |\eta| < 1.9 \): require at least 6 TRT hits, including TRT outliers, with outlier fraction under 90%;
  
  - Otherwise if \( |\eta| \leq 0.1 \) or \( |\eta| \geq 1.9 \): only if at least 6 TRT hits are observed, including TRT outliers, require the outlier fraction to be under 90%.

- Even more importantly, each muon candidate must pass strict requirements for Muon Spectrometer hits\(^2\):
  
  - To ensure that the momentum measurement is reliable, each muon track must have:
    
    * At least 3 hits in each of the BI, BM and BO MDT precision layers, or
    
    * At least 3 hits in each of the EI, EE and EM MDT precision layers, or
    
    * At least 3 hits in each of the EI, EM and EO MDT precision layers, or
    
    * At least 3 hits in each of the EM and EO MDT precision layers, along with at least 2 unspoiled CSC hits (unspoiled hits are the ones for which a precise \( \eta \) measurement is available);
    
  - As well, no hit is allowed in the MDT chambers not connected to the optical alignment system, namely the BEE, BIS7 and BIS8 chambers.

\(^2\) using the naming convention for MS chambers explained in Section 4.5.6
– For a correct estimate of the magnetic field along the track, it is necessary to require at least one hit measuring the $\phi$ coordinate, in two different layers of the RPC, TGC or CSC.

- To ensure that the momentum measurements are consistent between the different detector subsystems, for each muon candidate the difference between the standalone momentum measurements from the ID and MS must not exceed 5 times the sum in quadrature of the standalone uncertainties from each subsystem.

- In order to confirm that each muon candidate comes directly from the collision point, muon tracks are required to be close to the primary vertex: within 0.2 mm in the transverse direction ($d_0$) and within 1.0 mm in the longitudinal direction ($z_0$).

- The background due to the multi-jet background, discussed in Section 6.2.2, is significantly reduced by requiring each muon candidate to be isolated: specifically, the sum of the $p_T$ of all tracks with $p_T > 1$ GeV in a cone with radius $\Delta R = 0.3$ around the muon track must not amount to more than 5% of the muon $p_T$.

The selection used in 2011 at $\sqrt{s} = 7$ TeV is in general more restrictive than the one described above. First, muon candidates with hits in both the Barrel and one of the Endcaps could not be accepted in 2011, because the track reconstruction software did not take into account potential global mis-alignments of the Barrel with respect to the Endcaps. Following improvements in the track reconstruction software, this uncertainty is correctly propagated for the analysis at $\sqrt{s} = 8$ TeV. The algorithm now operates as follows: it is first determined whether a given muon traverses predominantly Barrel or Endcap chambers of the Muon Spectrometer, based on the number of hits. Then, hits from the corresponding region (Barrel or Endcap) are considered as usual in the track fits, while the error on the position hits from the other region is inflated by $\pm 7$ mm, corresponding to the global alignment uncertainty. In effect, this algorithm therefore discards hits from Endcap chambers when reconstructing muons from the Barrel, and similarly discards hits from Barrel chambers when reconstructing muons from the Endcap.

A second, major improvement for the analysis at $\sqrt{s} = 8$ TeV is the inclusion of muon candidates with hits in EE chambers. In 2011, most EE chambers were not installed, and the ones already installed had yet to be aligned and commissioned for use in analyses with high-momentum muons: muon tracks passing through these chambers were therefore vetoed by this analysis at the time. After the installation and alignment of many EE chambers in the winter shutdown between 2011 and 2012, use of these chambers was approved following dedicated studies of their resolution.

The ID hit requirements are a little tighter for the analysis in 2011: at least 2 Pixel hits are required instead of at least 1 (including Pixel dead sensors crossed), and at least 6 SCT hits are required instead of at least 5 (including SCT dead sensors crossed). Finally, at least 3 CSC hits in total are required, instead of at least 2 unspoiled hits.
5.3.2 Selection of 2-Station Muons

Additional gains in acceptance are attained by selecting the best muon candidates which have two track segments in the MS instead of three. While for 3-station muons, the main constraint on the track curvature comes from the sagitta measurement in the MS, for 2-station muons this information is unavailable. In this case the momentum measurement in the MS has to come from the angular difference \( \Delta \theta_{\text{seg}} \) between the two measured segments, in addition to the magnetic field integral along the track. The result of the full track fit can be approximated as:

\[
p = \frac{K}{\Delta \theta_{\text{seg}}} \int B \cdot dl \equiv \frac{K}{\Delta \theta_{\text{seg}}} B_{\text{int}}
\]

where \( K \) is a constant.

For this reason, 2-station muons are expected to exhibit a worse momentum resolution than 3-station muons, in addition to being more prone to large mis-measurement in cases where one of the MDT chambers is mis-aligned. An example of such a situation is illustrated in Figure 5.8: a position mis-alignment in one of the MDT chambers can cause the curvature of a track having only two MDT sections to be greatly underestimated, thereby resulting in a fake high-momentum muon. In other words, while for 3-station muons the momentum resolution is dominated by the sagitta bias, for 2-station muons the primary cause of resolution degradation is the error on the segment angle reconstruction in the bending direction.

Dedicated studies are therefore carried out to identify which 2-station muons pass through MS towers where the alignment precision is sufficient to guarantee a reliable momentum measurement. This is first attempted in the Barrel for 2-station muons with hits in the BI and BO chambers, as these muons have the best potential in terms of both acceptance gain and resolution.

**Study at \( \sqrt{s} = 7 \text{ TeV} \)**

As a first step, the angular resolution of segments in data, which by Equation 5.5 is translated into the momentum resolution of 2-station muons, is quantified in data using 3-station muons in each individual MS tower. Figure 5.9 demonstrates that the relationship between \( q/p \) and \( \Delta \theta_{\text{seg}}/(KB_{\text{int}}) \) is linear\(^3\), justifying this approach.

Since the resolution of the momentum measurement using the information from all three stations is negligible compared to that coming from the segment angular difference, the former is taken as the reference to which the latter is compared. The momentum resolution of 2-station muons can then be obtained from the standard deviation of the quantity \( p \cdot |\Delta \theta_{\text{seg}}|/(KB_{\text{int}}) \) from neighbouring 3-station muons. This is evaluated in each MS tower using Gaussian fits to the distribution. After an initial fit over the range \([0..2]\) to initialize values for the mean \( \mu \) and standard deviation \( \sigma \), the fit

\(^3\) The value \( K = 0.3 \) is used in this section for presentation purposes; it has no influence on the results.
**Figure 5.8:** Example of a potential track curvature mis-measurement for a 2-station muon where one of the inner MDT chambers is out of alignment. The actual muon path, that would correspond to the measurement using well-aligned BI and BO chambers, is represented with a continuous line. An error $\delta z$ in the position alignment of one of the chambers can cause the reconstructed track, represented by the dashed line, to be straighter. Figure credit: P.-F. Giraud.

**Figure 5.9:** Profile histograms demonstrating the linear relationship between $q/p$ and $\Delta\theta_{seg}/B_{int}$ for muon candidates in large (left) and small (right) $\phi$-sectors of the MS.

is repeated with a range restricted to $\pm 1.75\sigma$ of the previous fit until convergence. The momentum resolution is then given by $\sigma/\mu$.

Most towers have acceptable angular resolution: two typical fit results are shown in Figure 5.10. By contrast, a small number of towers are found where the alignment precision is insufficient for a reliable 2-station momentum measurement: examples of these are shown in Figure 5.11.
Figure 5.10: Example fit results for MS towers with good angular resolution in data at $\sqrt{s} = 7$ TeV, from a large $\phi$-sector (left) and from a small $\phi$-sector (right). These results are typical of most towers. The Gaussian fits are for illustration only.

Figure 5.11: Example fit results for MS towers with known poor angular resolution in data at $\sqrt{s} = 7$ TeV. The Gaussian fits are for illustration only; the distributions are clearly non-Gaussian, resulting in poor fits. These towers are vetoed in the analysis.
The selected 2-station muons must pass the same selection as described in the previous section for 3-station muons, with the exception that the following criteria for Muon Spectrometer hits are required instead:

- At least 5 hits in both the BI and BO MDT precision layers;
- At least one RPC hit to measure the $\phi$ coordinate;
- No hit is allowed in the MDT chambers not connected to the optical alignment system, namely the BEE, BIS7 and BIS8 chambers.
- In order to veto the towers with known poor angular resolution, no MDT hit:
  - in large (i.e. odd-numbered) $\phi$-sectors with $|\eta| > 0.85$,
  - in $\phi$-sector 2 with $|\eta| > 0.85$,
  - or in $\phi$-sector 13 with $0.00 < \eta < 0.65$;

In addition to the specific requirements on the $\eta$ coordinate used to veto muon candidates passing through the vetoed towers, a general requirement of $|\eta| < 1.00$ is imposed to make sure to select muons from Barrel chambers only. Additionally, the difference between the standalone momentum measurements from the ID and MS must not exceed 3 times the sum in quadrature of the standalone uncertainties. This requirement is tighter with respect to that for 3-station muons, because 2-station muons rely more heavily on the ID track measurement.

With this selection in hand, it is then possible to study the resolution of 2-station muons as a function of momentum. Figure 5.12 shows fit results for selected 2-station muons from data in the momentum bin 200–300 GeV, taken as an example. The momentum resolution is obtained for each momentum bin in the same way as for individual towers. The result is shown for muon candidates from data and simulated muons in Figures 5.13 and 5.14 respectively.

The resolution parameters specific to high-momentum 2-station muons $p_{MS}^{\star,2,MC}$ are then derived respectively for data and simulation using fits with the functional form of Equation 5.1. The results are shown in Table 5.5, along with the corresponding smearing parameters given by $\sigma_{2}^{MS} = p_{2,\text{data}}^{\star,2,MC} = p_{2,\text{MC}}^{\star,2,MC}$.

As a cross-check, another way to obtain the momentum resolution of simulated 2-station muons is to compare their momentum values as measured by the virtual ATLAS detector to the real values (known as “MC truth”) from the event generator. The resolution is then given by the standard deviation of the quantity $p/p_{\text{truth}} - 1$. Figure 5.15 shows the results as a function of momentum. Using a fit with the same parametrization, the resolution parameter $p_{2,\text{MC}}^{\star,2}$ is measured to be $0.41 \pm 0.02$ for both the large and small $\phi$-sectors, in agreement with the segment angular difference method.

In addition to the 2-station muons with hits in the BI and BO chambers studied in this section, accepting additional 2-station muons with hits in the BI and BM chambers or in the BM and BO
chambers was briefly considered. Because for these muons the distance between the two MS stations is half the distance between the Inner and Outer stations, angular differences are in general only half as large: as a result the angular resolution of segments fails to provide satisfactory momentum resolution for 2-station muons at very high momentum. Specifically, preliminary values for $P_{2\text{MS}}^{\ast}$ above 100%/TeV were observed for 2-station muons with hits in the BI and BM chambers or in the BM and BO chambers. Similar conclusions were reached when considering 2-station muons in the Endcaps with hits in the EI and EM chambers, in towers where the EE chambers were not yet installed.
Figure 5.14: Momentum resolution of 2-station muons in simulation at $\sqrt{s} = 7$ TeV. Here the resolution of simulated muons is obtained using the same method as for data.

Figure 5.15: Momentum resolution of 2-station muons in simulation at $\sqrt{s} = 7$ TeV from comparing the refit 2-station momentum to MC truth.
Table 5.5: $p_{MS}^{2\ast}$ parameters measured in data and simulation at $\sqrt{s} = 7$ TeV and corresponding smearing parameters $S_{MS}^{2\ast} = p_{MS,\text{data}}^{2\ast} \ominus p_{MS,\text{MC}}^{2\ast}$.

<table>
<thead>
<tr>
<th>Sectors</th>
<th>$p_{2,\text{data}}^{MS\ast}$ [TeV$^{-1}$]</th>
<th>$p_{2,\text{MC}}^{MS\ast}$ [TeV$^{-1}$]</th>
<th>$S_{2}^{MS\ast}$ [TeV$^{-1}$]</th>
</tr>
</thead>
<tbody>
<tr>
<td>Small</td>
<td>$0.55 \pm 0.02$</td>
<td>$0.45 \pm 0.02$</td>
<td>$0.32 \pm 0.04$</td>
</tr>
<tr>
<td>Large</td>
<td>$0.54 \pm 0.02$</td>
<td>$0.41 \pm 0.02$</td>
<td>$0.35 \pm 0.04$</td>
</tr>
</tbody>
</table>

Table 5.6: $p_{2}^{MS\ast}$ parameters measured in data and simulation at $\sqrt{s} = 8$ TeV.

<table>
<thead>
<tr>
<th>Sectors</th>
<th>$p_{2,\text{data}}^{MS\ast}$ [TeV$^{-1}$]</th>
<th>$p_{2,\text{MC}}^{MS\ast}$ [TeV$^{-1}$]</th>
</tr>
</thead>
<tbody>
<tr>
<td>Small</td>
<td>$0.46 \pm 0.02$</td>
<td>$0.54 \pm 0.02$</td>
</tr>
<tr>
<td>Large</td>
<td>$0.42 \pm 0.02$</td>
<td>$0.48 \pm 0.02$</td>
</tr>
</tbody>
</table>

Results of the study at $\sqrt{s} = 8$ TeV

The same study is repeated for the analysis at $\sqrt{s} = 8$ TeV. The requirements on 2-station muons are re-assessed and turn out to be the same as for the dataset collected in 2011, with the exception of the requirements on the $\eta$ coordinate and vetoed chambers. The general requirement on $\eta$ for 2-station muons is relaxed to $|\eta| < 1.05$, and following a repetition of the tower-by-tower study, the vetoed chambers are located:

- in $\phi$-sectors 4 or 6 with $|\eta| > 0.85$,
- in $\phi$-sector 9 with $0.20 < |\eta| < 0.35$,
- and in $\phi$-sector 13 with $0.00 < \eta < 0.20$.

Therefore the alignment of many towers vetoed in the 2011 dataset has improved in the 2012 dataset.

The measured resolution parameters for these muon candidates are shown in Table 5.6. Like for the 3-station muons in this dataset, the simulated resolution parameter $p_{2}^{MS\ast}$ is observed to be better in data than in simulations, and therefore no additional smearing is applied. Since the differences in resolution between the simulation and data are small for 3-station muons, of which the selected sample is mostly composed, the differences in resolution for 2-station muons are not expected to impact the sensitivity of the search.
Chapter 6

Event Selection and Comparison of Data with Standard Model Expectations

This chapter first describes the criteria used to select the events considered in the search, in Section 6.1. Then, Section 6.2 explains the techniques used to evaluate the backgrounds to the search and the expected signal contributions from the new physics models under consideration. Finally, Section 6.3 compares the observed data to these predictions.

6.1 Event Selection

To be considered in the search for high-mass resonances decaying to muon pairs, collision events are required to pass the following requirements:

- Candidate events have to satisfy data-quality requirements, as discussed in Section 4.8.3.
- Candidate events must have triggered the detector with at least one of the following single-muon triggers, requiring:
  - For events collected in 2011 at $\sqrt{s} = 7$ TeV:
    * At least one combined muon with $p_T > 22$ GeV, or
    * At least one standalone muon in the Barrel with $p_T > 40$ GeV.
  - For events collected in 2012 at $\sqrt{s} = 8$ TeV:
    * At least one combined muon with $p_T > 24$ GeV and an isolation requirement (see below), or
    * At least one combined muon with $p_T > 36$ GeV and no isolation requirement.
- The primary vertex of the event, defined as the one with the largest $\sum p_T^2$, where the sum is over all tracks in the ID with $p_T > 0.4$ GeV, must have at least 3 tracks, and be located within 20 cm of the centre of the detector along the beam axis.
• Candidate events must have at least two high-momentum muons passing the selection detailed in Section 5.3.

The standalone muon trigger chain is used to recover inefficiencies in the combined muon trigger chains for data collected at $\sqrt{s} = 7$ TeV. For data collected at $\sqrt{s} = 8$ TeV, the isolation requirement used in the trigger with a lower $p_T$ threshold is that the sum of the $p_T$ of all tracks with $p_T > 1$ GeV in a cone with radius $\Delta R = 0.2$ around the muon track must not amount to more than 12% of the muon $p_T$. This requirement at trigger-level is significantly looser than the isolation requirement used in the analysis. Nevertheless, a non-isolated muon trigger chain is used to recover possible inefficiencies, due to differences between the tracks used in the calculation of the isolation variable at trigger-level as compared with the tracks used for this purpose in the final recorded event.

In each candidate event, high-momentum muons are used to build one opposite-sign muon pair. Following the classification explained in Section 5.3, the selected muon candidates can be 3-station or 2-station muons. First, if two opposite-sign 3-station muon candidates are found, they are used to make the pair: the pair is then said to pass the “primary dimuon selection”. If not, the pair can be built using one 3-station muon and one 2-station muon, in which case it is said to pass the “secondary dimuon selection”. If more than one pair passing the primary selection is found, the one with the highest transverse momentum scalar sum is selected; similarly, if the event has no muon pair passing the primary selection but more than one passing the secondary selection, the muon pair with the highest transverse momentum scalar sum is selected. Finally, events where the electric charge of all muons is of the same sign are discarded.

Figure 6.1 shows the signal acceptance times efficiency for a $Z'_\text{SSM}$ boson as a function of $M_{Z'}$. The values at $M_{Z'} = 2.5$ TeV for the analysis on the dataset collected in 2012 are 40.0% for the primary dimuon selection and 3.2% for the secondary dimuon selection; these numbers are respectively 36.2% and 2.9% for the analysis using data collected in 2011. The reasons for the improvement between 2011 and 2012 are explained in Section 5.3.

Figures 6.2 and 6.3 show the event yield per pb$^{-1}$ per run for the full event selection described above. The yields are fairly constant as a function of time. For the analysis at $\sqrt{s} = 7$ TeV, two slight efficiency losses are noticeable, due to the following causes:

• A tighter Level-1 muon trigger, necessary to handle a higher instantaneous luminosity, was used from the start of data period J (run 186516);

• A timing problem in the RPC affects the trigger efficiency for runs 189205–189610 in data period L.

The observed event with the highest dimuon invariant mass is displayed in Figure 6.4. The event passes the primary dimuon selection, with the two 3-station muons having transverse momenta of 652 GeV and 646 GeV respectively. The invariant mass of the muon pair is 1.84 TeV.
Figure 6.1: Signal acceptance times efficiency for a $Z_{SSM}'$ boson as a function of $M_{Z'}$. The region $M_{Z'} > 3$ TeV was not considered for the analysis on the 2011 dataset. The width of the lines is representative of the uncertainty.
Figure 6.2: Yield per pb$^{-1}$ for each run, for the primary (top) and secondary (bottom) dimuon selection at $\sqrt{s} = 7$ TeV.
Figure 6.3: Yield per pb$^{-1}$ for each run, for the primary (top) and secondary (bottom) dimuon selection at $\sqrt{s} = 8$ TeV.
Figure 6.4: Event with the highest dimuon invariant mass observed in data collected by the ATLAS experiment in Run-I of the LHC. The muon candidates have transverse momenta of 652 GeV and 646 GeV respectively, and the dimuon invariant mass is 1.84 TeV.
6.2 Background and Signal Expectation

The main background contributions in this analysis are evaluated using simulated samples, described in Section 6.2.1. Contributions from multi-jet collision events and cosmic rays are evaluated using data-driven techniques, and turn out to be negligible: these techniques are described in detail in Section 6.2.2. Section 6.2.3 details how signal templates are obtained.

6.2.1 Simulated Samples

As already mentioned in Section 2.2.2, since protons are not fundamental particles, being made of quarks and gluons (collectively partons), in a typical proton-proton collision event one parton from each incoming proton interacts in the main process of the event. Softer, simultaneous interactions constitute the underlying event, while partons that interact minimally with the proton bunch traveling in the opposite direction are called beam remnants.

Figure 6.5: Sketch of a proton-proton collision at high energy. Figure credit: F. Siegert.

Figure 6.5 displays a sketch of such a high-energy proton-proton collision event. The event illustrated is a fully hadronic process: both the initial and the final states of the main process are constituted of quarks and gluons, which along with any emitted strongly coupled radiation undergo parton showering, fragmentation and hadronization before reaching the detector. The resulting object
is called a hadronic jet. In comparison, muons are relatively easier to model, because they only emit electroweak radiation as they traverse the detector. It is nevertheless crucial that event simulations used to calculate signal and background expectations take into account all of the effects described above if they are to describe collision data accurately.

**Event yield predictions using Monte Carlo simulations**

Considering an hypothesis $H$ in a region of phase space, the expected number of collision events selected by the analysis is given by

$$\mu = L_{\text{int}} \sum_{i \in H} (\sigma B A \varepsilon)_i$$

(6.1)

where $L_{\text{int}}$ is the integrated luminosity, $\sigma_i$ is the cross section and $B_i$ the branching fraction of a given physical process $i$ into final states of interest, $A_i$ is the acceptance of the detector for this process and $\varepsilon_i$ is the efficiency for the accepted events to pass the analysis selection criteria. If $H = H_0$ is the null hypothesis, then $\mu$ represents the background expectation; on the other hand if $H = H_{Z'}$ includes a non-zero signal strength from a given $Z'$ boson, then $\mu$ represents the expectation from the sum of background and $Z'$ signal.

In practice, the expected value $\mu$ is obtained using Monte Carlo simulation techniques, by generating events for the physical processes of interest and simulating their interaction with a virtual representation of the ATLAS detector. The quantity $(A \varepsilon)_i$ is then the fraction of simulated events that are accepted by the virtual detector and pass the analysis selection criteria, over the total number of generated events:

$$A \varepsilon = \frac{N_{\text{pass}}}{N_{\text{gen}}}$$

(6.2)

Corrections described in the remainder of this section are applied to both $N_{\text{pass}}$ and $N_{\text{gen}}$. The passing events are then scaled by an additional weight $w_i = L_{\text{int}}(\sigma B / N_{\text{gen}})_i$ such that their weighted sum gives the expectation:

$$\mu = \sum_{i \in H} (w N_{\text{pass}})_i = \sum_{i \in H} \left( L_{\text{int}} \sigma B \frac{N_{\text{pass}}}{N_{\text{gen}}} \right)_i$$

(6.3)

This expression can be used to find the expected number of events in any region of phase space: these regions are typically represented by histogram bins in each variable of interest.

**Analysis at $\sqrt{s} = 7$ TeV**

Table 6.1 indicates the event generators used for the hard process and parton shower, as well as the PDFs used to generate simulated samples for the analysis at $\sqrt{s} = 7$ TeV. Such samples are used to obtain $Z'$ signal templates and to evaluate the background yields from $Z/\gamma'$, $WW$, $WZ$,
Table 6.1: Summary of the simulated samples for the analysis using the full dataset at $\sqrt{s} = 7$ TeV. The re-weighting of $Z'$ samples is discussed in Section 6.2.3.

<table>
<thead>
<tr>
<th>Process</th>
<th>Main generator</th>
<th>Parton shower</th>
<th>PDFs</th>
</tr>
</thead>
<tbody>
<tr>
<td>$t\bar{t}, Wt$</td>
<td>MC@NLO 4.01 [101]</td>
<td>HERWIG 6.510</td>
<td>CTEQ66 [137]</td>
</tr>
<tr>
<td>$Z'$</td>
<td>PYTHIA 6.421 (re-weighted)</td>
<td>PYTHIA 6.421</td>
<td>MRST2007LO**</td>
</tr>
</tbody>
</table>

ZZ, $t\bar{t}$ and $Wt$ production in final states with at least two muons of opposite charge. With the exception of the samples involving top quarks, which are generated using matrix elements at NLO, all other samples are generated using LO event generators, making use of the modified LO PDF set MRST2007LO** [156, 157]. This modified PDF set was developed in an attempt to emulate the calculation results from NLO generators when using LO generators. For all samples, the program PHOTOS [110] is used to simulate final-state photon radiation. The interactions of the generated particles with the ATLAS detector are then simulated using GEANT4 [104], and the events are fully reconstructed using the ATLAS reconstruction software [15].

The dominant, irreducible background to this search is $Z/\gamma^* \rightarrow \mu^+\mu^-$. To obtain the state-of-the-art prediction for the yields due to this process as a function of dimuon invariant mass, the $Z/\gamma^*$ cross section is calculated using PHOZPR [114] at NNLO in perturbative QCD, with MSTW2008NNLO PDFs [132]. The simulated $Z/\gamma^*$ samples are then re-weighted event-by-event using the ratio of this NNLO cross section to the prediction from the LO simulation: this correction factor is shown in Figure 6.6. Multiplicative correction factors such as this one are commonly called “K-factors”.

Another event-by-event weight correction, displayed as a function of $m_{\mu^+\mu^-}$ in Figure 6.7, is also applied to account for NLO electroweak corrections due to virtual heavy gauge boson loops, photon-induced processes ($\gamma\gamma \rightarrow \mu^+\mu^-$) and real radiation of $W$ and $Z$ bosons from the final-state muons. The first two corrections are calculated at LO using HORACE 3.1 [64] with MRST2004QED PDFs [131]. Real radiation is estimated based on the results of Ref. [48] to contribute an enhancement of about 2% / TeV. The inclusive $Z/\gamma^*$ cross section with the above higher-order corrections, for $m_{\mu^+\mu^-} > 60$ GeV, is $\sigma_{Z/\gamma^*} = 989 \pm 49$ pb.

Additional corrections not changing the total background normalization are applied to the weights of simulated events to account for known discrepancies with respect to data. To correct for differences between the amount of pileup events in the simulated samples with respect to the observed distribution shown in Figure 4.6, a correction is applied as a function of the number of interactions per proton bunch crossing. The corrections to simulated muons due to the detector efficiency and resolution described in Section 5.2 are applied as well.

In addition, a significant difference was found in the $p_T$ spectrum of $Z$ bosons in simulated
Figure 6.6: Perturbative QCD correction factor for $Z/\gamma^* \rightarrow \mu^+\mu^-$ production as function of $m_{\mu^+\mu^-}$ at $\sqrt{s} = 7$ TeV. The purple dash-dotted line, representing the correction from the Leading Order calculation with MRST2007LO$^*$ to the Next-to-Next-to-Leading Order calculation with MSTW2008NNLO, is used in the analysis at $\sqrt{s} = 7$ TeV. Figure credit: T. Nunnemann.

Figure 6.7: Electroweak and photon-induced correction factor for $Z/\gamma^* \rightarrow \mu^+\mu^-$ production as function of $m_{\mu^+\mu^-}$ at $\sqrt{s} = 7$ TeV, taking into account corrections due to virtual $W$ and $Z$ loops and real final-state radiation of $W$ and $Z$ bosons. A piecewise polynomial fit to the correction is shown in red. Figure credit: T. Nunnemann.
samples produced at $\sqrt{s} = 7$ TeV in 2011, compared to similar samples generated at the same centre-of-mass energy in 2010. This difference can be seen in Figure 6.8. While the earlier production is found to describe collision data accurately, the later one does not. This difference was traced back to a re-tuning of event generator settings. A correction is therefore applied event-by-event as a function of the generated $Z p_T$ and $m_{\mu^+\mu^-}$ to recover the spectrum from 2010 simulation.

Figure 6.8: Generated $Z p_T$ spectra near the $Z$ resonance region (75 GeV < $m_{\mu^+\mu^-}$ < 120 GeV, in black) and at higher invariant masses (600 GeV < $m_{\mu^+\mu^-}$ < 800 GeV, in red). Samples generated in 2010 and 2011 are respectively shown with closed and open markers. Figure credit: J. Kretzschmar.

Top-quark backgrounds, namely $t\bar{t}$ and $Wt$ production, also contribute opposite-sign pairs of high-momentum muons when both $W$ bosons in the event decay to $\mu\nu$. The $t\bar{t}$ cross section is scaled to an approximate-NNLO prediction of $\sigma_{t\bar{t}} = 160^{+11}_{-15}$ pb [125, 135], and the single-top $Wt$ cross section used is $\sigma_{Wt} = 14.4 \pm 1.0$ pb. Diboson backgrounds evaluated from $WW$, $WZ$ and $ZZ$ simulated processes are similarly scaled to NLO cross sections, calculated using MCFM [63] with an uncertainty of 5%. The values used are $\sigma_{WW} = 45 \pm 2$ pb, $\sigma_{WZ} = 18.0 \pm 0.9$ pb and $\sigma_{ZZ} = 6.0 \pm 0.3$ pb.

Because both the top-quark and diboson simulated background samples lack sufficient statistics at very high dimuon invariant masses, an extrapolation is performed by fitting the background predictions, with all corrections applied, to the functional form

$$N(m) = p_1 \cdot m^{p_2} \cdot \log m + p_3$$

where the $p_i$ are fit parameters and $m$ is the invariant mass $m_{\mu^+\mu^-}$. For top-quark backgrounds, the fit region is 200 GeV < $m_{\mu^+\mu^-}$ < 800 GeV, while for diboson backgrounds it is 450 GeV < $m_{\mu^+\mu^-}$ <
Table 6.2: Summary of the simulated samples for the analysis using the full dataset at √s = 8 TeV. The re-weighting of Z' samples is discussed in Section 6.2.3.

<table>
<thead>
<tr>
<th>Process</th>
<th>Main generator</th>
<th>Parton shower</th>
<th>PDFs</th>
</tr>
</thead>
<tbody>
<tr>
<td>Z/γ*</td>
<td>POWHEG [3]</td>
<td>PYTHIA 8.162 [159]</td>
<td>CT10 [121]</td>
</tr>
<tr>
<td>WW, WZ, ZZ</td>
<td>HERWIG++ 2.5.2 [82, 83]</td>
<td>HERWIG 6.520</td>
<td>CTEQ6L1</td>
</tr>
<tr>
<td>t̅t, Wt</td>
<td>MC@NLO 4.06 [101]</td>
<td>HERWIG 6.520</td>
<td>CT10</td>
</tr>
<tr>
<td>Z'</td>
<td>PYTHIA 8.165 (re-weighted)</td>
<td>PYTHIA 8.165</td>
<td>MSTW2008LO [132]</td>
</tr>
</tbody>
</table>

1450 GeV. Above these ranges in invariant mass, where the simulated statistics are insufficient, the background predictions are taken from Equation 6.4 using the parameters from the corresponding fits. Two sources of systematic uncertainties on this extrapolation are considered. First, the fit range is varied by 5 steps of 10 GeV about the starting value and 5 steps of 20 GeV before the ending value, resulting in 25 different fits, and the maximum discrepancy with respect to the nominal fit is taken. Second, a different functional form is tried: \( N(m) = p_1 \cdot (m + p_2)^{p_3} \). The sum in quadrature of the effects from these two variations is taken as the uncertainty on the estimated background yield where it is taken from the fit.

In the analysis using the first 1.21 fb^{-1} of data collected at √s = 7 TeV, the background from W + jets processes, where one muon comes from the W → µν decay in the hard scatter and the other from a meson decay in one of the jets, was also estimated. Simulated samples were generated using ALPGEN [128] interfaced with CTEQ6L1 PDFs [146], with HERWIG to simulate the parton showers and underlying event and JIMMY 4.31 [59] for multiple parton interactions. This background turns out to be negligible, as can be seen in Figure 6.12, and was therefore not considered in subsequent iterations of the analysis.

Finally, the sum of all backgrounds is normalized in the Z peak region, corresponding to \( m_{µ^+µ^-} \) between 70 and 110 GeV. The number of events passing the full selection found in this normalization region is 985,180 and the corresponding scale factor is 1.01, with an uncertainty of 5% dominated by that due to the theoretical Z cross section. This scale factor value is consistent with unity, and its central value lies within the 1.8% uncertainty on the integrated luminosity measurement.

**Analysis at √s = 8 TeV**

For the analysis using data collected at √s = 8 TeV, background estimates are obtained following the same general strategy, using updated versions of the same event generators except for the Z/γ* background, where POWHEG is used instead of PYTHIA for the hard process. The event generators used for the hard process and parton shower, as well as the PDFs used, are indicated in Table 6.2.

Here the state-of-the-art prediction for the Z/γ* background is calculated at NNLO in perturbative QCD with NLO electroweak corrections using FEWZ [126, 134] with MSTW2008NNLO PDFs. The resulting K-factors applied to samples generated with POWHEG and PYTHIA are shown.
This calculation also includes photon-induced contributions, which are estimated at LO with MRST2004QED PDFs, as well as real final-state radiation of $W$ and $Z$ bosons estimated using MADGRAPH 5 [5] following the prescription outlined in Ref. [48]. The electroweak and photon-induced corrections were verified by SANC [42, 53]. The $Z/\gamma^*$ cross section thus calculated, integrated for $m_{\mu^+\mu^-} > 60 \text{ GeV}$ is $\sigma_{Z/\gamma^*} = 1147 \pm 50 \text{ pb}$.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{Figure6.9.png}
\caption{Perturbative QCD, electroweak and photon-induced correction factors for $Z/\gamma^* \rightarrow \mu^+\mu^-$ production as function of $m_{\mu^+\mu^-}$ at $\sqrt{s} = 8 \text{ TeV}$. The electroweak correction factors take into account corrections due to virtual $W$ and $Z$ loops and real final-state radiation of $W$ and $Z$ bosons. Calculation by U. Klein.}
\end{figure}

In a manner identical as explained above for the analysis at $\sqrt{s} = 7 \text{ TeV}$, the appropriate corrections to the number of interactions per proton bunch crossing as well as to the detector efficiency and resolution for reconstructing muons are also applied. On the other hand, no correction to the simulated $Z p_T$ is necessary, following improvements in the Monte Carlo generator tune.

Top-quark backgrounds are scaled to the cross section $\sigma_{t\bar{t}} = 253^{+13}_{-15} \text{ pb}$, obtained from a calculation performed at NNLO in perturbative QCD including the re-summation of next-to-next-to-leading logarithmic soft gluon terms with Top++ 2.0 [44, 61, 85-88], assuming a top quark mass of 172.5 GeV. The uncertainties on this prediction include PDF and $\alpha_S$ uncertainties calculated using the PDF4LHC prescription [55] with the NNLO MSTW2008 [132, 133], CT10 [103, 121] and NNPDF2.3 [41] PDF error sets, QCD scale uncertainties and the effect of varying the top quark mass by $\pm 1 \text{ GeV}$. The single-top $Wt$ cross section is taken to be $\sigma_{Wt} = 22.4 \pm 1.5 \text{ pb}$ [119].
NLO cross section calculation for diboson processes is repeated at $\sqrt{s} = 8$ TeV with MCFM, and the values obtained are $\sigma_{WW} = 57 \pm 3$ pb, $\sigma_{WZ} = 21.5 \pm 1.1$ pb and $\sigma_{ZZ} = 7.4 \pm 0.4$ pb.

For this version of the analysis, the diboson samples have sufficient statistics over the full $m_{\mu^+\mu^-}$ spectrum. Top background contributions are extrapolated using a fit like for the analysis at $\sqrt{s} = 7$ TeV, using the functional form of Equation 6.4. The uncertainties are derived in the same way, by varying the fit range and functional form.

The normalization region is taken to be $80$ GeV $< m_{\mu^+\mu^-} < 110$ GeV. The number of events passing the full selection found in this normalization region is 5,075,739 and the corresponding scale factor is 0.98, with a 4% theoretical uncertainty from the Z cross section. Here too, the scale factor value is consistent with unity, and its central value lies within the preliminary 3.6% uncertainty on the integrated luminosity measurement.

6.2.2 Data-Driven Background Estimates

This section describes the data-driven background estimates of contributions from multi-jet processes and cosmic rays, using partial datasets collected at $\sqrt{s} = 7$ TeV. Both of these backgrounds were estimated to be negligible then, and re-evaluating them using the complete dataset is not necessary.

Multi-jets

In hadron colliders, fully hadronic processes have cross sections ranking orders of magnitude above the cross sections for hard-scattering processes with leptons in the final state; for illustration see Figure 2.7 in Section 2.2.2. Following hadronization, jets originally from quarks and gluons can have component particles which decay into muons, e.g. $\pi^\pm \rightarrow \mu^\pm \nu$ or $K^\pm \rightarrow \mu^\pm \nu$. While these muons from jets generally fail the tight isolation and impact parameter requirements used in this analysis, they constitute a potential background to the search in events where two of them pass all selection requirements.

Using the first 1.21 fb$^{-1}$ of data collected at $\sqrt{s} = 7$ TeV in 2011, a pure sample of muon pairs where both muons come from jets is obtained by reverting the requirement on the isolation variable $R_{iso} = \Sigma p_T^{tag}(\Delta R < 0.3)/p_T^{\mu}$, to $0.1 < R_{iso} < 1.0$ instead of $R_{iso} < 0.05$ in the nominal selection. The distribution in the muon isolation variable immediately before the isolation requirement is applied is shown on Figure 6.10 and the invariant mass distribution of events where two muon candidates pass the reversed isolation requirement along with the rest of the event selection is displayed in Figure 6.11. In both these histograms, the prediction from a simulated multi-jet background sample generated using PYTHIA is shown.

The shape for the QCD multi-jet background is then taken from this Control Region (CR), and
Figure 6.10: Muon track-based isolation distribution, immediately before the isolation cut.

The QCD multi-jet background (yellow) is taken from PYTHIA MC simulation scaled by a factor 0.5. Extrapolated to the Signal Region (SR) by applying a scaling factor as follows:

\[
\frac{N_{\text{exp},\text{SR}}}{N_{\text{exp},\text{CR}}} = \left( \frac{P_{\text{MC}}(R_{\text{iso}} < 0.05)}{P_{\text{MC}}(0.1 < R_{\text{iso}} < 1.0)} \right)^2
\]

\[
= (0.024 \pm 0.012)^2
\]

\[
= (5.8 \pm 5.8) \times 10^{-4}
\]

where \(P_{\text{MC}}(R_{\text{iso}})\) is the probability for a muon from multi-jet processes to fall within a certain range in the isolation variable, evaluated using PYTHIA simulation. The ratio has a statistical uncertainty of 50%, which propagates to a 100% uncertainty on the multi-jet estimate in the SR.

The resulting multi-jet background estimate can be seen in Figure 6.12, which shows the dimuon invariant mass distribution after final selection for the first 1.21 fb\(^{-1}\) of data collected in 2011. The multi-jet background is smaller than dominant backgrounds by many orders of magnitude, and therefore negligible in this analysis.
Figure 6.11: Dimuon invariant mass distribution from events with two muons passing the reversed isolation requirement along with the rest of the selection. The QCD multi-jet background (yellow) is taken from PYTHIA MC simulation scaled by a factor 0.5.

Figure 6.12: Dimuon invariant mass distribution after final selection for the first 1.21 fb$^{-1}$ of data collected in 2011. Here the QCD multi-jet background (yellow) corresponds to the data-driven estimate, i.e. the data shown in Figure 6.11 with extrapolation factor applied. The cosmic ray background is too small to be shown.
Cosmic rays

In addition to being produced in the LHC collisions, high-momentum muons can also come from cosmic rays. They are produced when high-energy protons from astronomical sources collide with molecules in the Earth’s atmosphere, each producing a cascade of hadronic particles. Among these particles produced are mesons such as pions and kaons, which often decay into muons.

Just like muons originating from the collider, high-momentum cosmic muons are minimum-ionizing particles, and have long lifetimes due to relativistic time dilation. They can therefore traverse the atmosphere and the shielding above the ATLAS detector, before going through the detector itself. In the event that a cosmic muon crosses all detector subsystems and coincides with a collision event at the interaction point before exiting under the detector, it can be mistakenly reconstructed as a prompt dimuon event. An example of such an event, observed during a cosmic-ray run in 2008, is displayed in Figure 6.13.

**Figure 6.13:** Cosmic muon event with hits in all Barrel detectors. Both solenoid and toroid magnets were operational during this run, taken in 2008 [33].
A study estimating the cosmic ray background contamination in the high-mass dimuon sample is conducted on the partial dataset consisting of the first 236 pb\(^{-1}\) of data collected in 2011. This corresponds to the first 1.5 months of data taking, out of 7.5 months that year. A sample of probable cosmic ray events is obtained from this collision data by reverting the impact parameter cuts used in the event selection, requiring 0.3 mm < \(d_0\) < 10 mm for the transverse impact parameter and 3.0 mm < \(z_0\) < 200 mm for the impact parameter along the beam axis (instead of \(d_0\) < 0.2 mm and \(z_0\) < 1.0 mm in the nominal selection). Loose upper cuts on the impact parameters are kept for the CR, in order to maintain the same efficiency for detecting muons in the CR as in the SR.

Two events are found to pass these reversed impact parameter cuts along with the rest of the event selection. This number in the CR (\(N_{\text{cosmics,CR}}\)) is extrapolated to the number of cosmic ray events expected in the SR (\(N_{\text{cosmics,SR}}\)). Assuming a constant cosmic ray distribution in spatial coordinates perpendicular to the vertical, an upper bound is obtained by scaling the number in the CR using the ratio of the dimensions of the SR and CR along the \(x\)- and \(z\)-axes. This constitutes an upper bound because considering any correction along the \(y\)-axis would reduce the estimate. This yields

\[
N_{\text{exp,SR}} < N_{\text{data,CR}} \cdot \frac{d_{0,\text{SR}}}{d_{0,\text{CR,outer}} - d_{0,\text{CR,inner}}} \cdot \frac{z_{0,\text{SR}}}{z_{0,\text{CR,outer}} - z_{0,\text{CR,inner}}} < (2.0 \pm 1.4) \cdot \frac{0.2 \text{ mm}}{10 \text{ mm} - 0.3 \text{ mm}} \cdot \frac{1.0 \text{ mm}}{200 \text{ mm} - 3.0 \text{ mm}} < (2.0 \pm 1.4) \cdot (2 \times 10^{-2}) \cdot (5 \times 10^{-3}) < (2.0 \pm 1.4) \times 10^{-4}
\]

Since the number of cosmic rays entering the detector scales with trigger live-time, the corresponding bound on the number of cosmic events passing all selection requirements for the full \(\sqrt{s} = 7\) TeV dataset is

\[
(2.0 \pm 1.4) \times 10^{-4} \cdot \frac{7.5 \text{ months}}{1.5 \text{ months}} = (10 \pm 7) \times 10^{-4}
\]

The cosmic ray contamination is therefore negligible in this analysis.

Since the trigger live-time and the selection are similar for the \(\sqrt{s} = 8\) TeV dataset, the cosmic ray contamination is expected to be negligible there as well.

### 6.2.3 Generation of Signal Templates

This search makes use of signal templates as a function of the dimuon invariant mass \(m_{\mu^+\mu^-}\), to interpret the significance of differences between observed data and the expected backgrounds. For each signal hypothesis, the combined signal and background templates are compared to the data.
using a likelihood fit: this is described in detail in Chapter 8. The main advantage of the template method over a simple counting method using an invariant mass window is that templates make use of more information by taking into account signal contributions across the whole mass spectrum, therefore increasing the signal acceptance and allowing a better sensitivity.

Nominally, a dedicated signal template is obtained for each \( Z' \) pole mass \( M_{Z'} \). In this approach, the relative width \( \Gamma_{Z'}/M_{Z'} \) of the resonance is fixed, corresponding to that of the specific \( Z' \) model under study: above the \( \bar{t}t \) production threshold, its value is 3.0\% for \( Z'_{\text{SSM}} \), 1.2\% for \( Z'_{\chi} \) and 0.5\% for \( Z'_{\psi} \). The iterations of the analysis at \( \sqrt{s} = 7 \) TeV and \( \sqrt{s} = 8 \) TeV use different techniques to generate these signal templates; both are described below. Moreover, this section also describes a third technique, used for generating signal templates as a function of the coupling \( \gamma' \) and mixing angle \( \theta_{\text{Min}} \) from the Minimal \( Z' \) Models described in Section 3.2.2. These templates have the additional advantage to fully take into account the interference between the \( Z/\gamma' \) and \( Z' \) processes.

In all cases, the same mass-dependent higher-order perturbative QCD corrections described in Section 6.2.1 are applied to \( Z' \) signal template shapes, under the assumption that differences between the initial-state quarks couplings to \( Z' \) bosons and their couplings to Z bosons do not significantly change these corrections. Higher-order electroweak corrections are not applied to the \( Z' \) signal, since as mentioned in Section 3.2.1, the unknown couplings between \( Z' \) and heavy gauge bosons are assumed to be zero. Corrections due to experimental effects are applied, namely to the muon efficiency and momentum resolution, as well as the number of interactions per bunch crossing.

**Nominal \( Z' \) signal templates**

For the nominal analysis at \( \sqrt{s} = 7 \) TeV, a sample of signal events is generated with a constant distribution as a function of \( m_{\mu^+\mu^-} \), by modifying the expression for the differential \( Z' \) cross section in PYTHIA at generator-level. This is accomplished first by removing the dependence on \( M_{Z'} \) and \( \Gamma_{Z'} \), by taking out a Breit-Wigner factor: this leaves a smoothly-falling event distribution as a function of \( m_{\mu^+\mu^-} \). An exponential factor is then used to make this spectrum constant as a function of \( m_{\mu^+\mu^-} \). Mathematically, the transformation is the following:

\[
\frac{d\sigma}{dm^2 dy} \rightarrow \frac{d\sigma}{dm^2 dy} \times ((m^2 - M_{Z'}^2)^2 + M_{Z'}^4 \Gamma_{Z'}^2) \times \exp(0.00195 \cdot m)
\]

(6.8)

where \( m \) is the dimuon invariant mass \( m_{\mu^+\mu^-} \). From this uniform distribution of events, templates can then be obtained at any \( Z' \) pole mass value using an event-by-event weight \( \mathcal{W} \), multiplying back a Breit-Wigner factor corresponding to the \( M_{Z'} \) and \( \Gamma_{Z'} \) values of interest and the reciprocal of the same exponential factor. Explicitly,

\[
\mathcal{W}(m) = \frac{\exp(-0.00195 \cdot m)}{(m^2 - M_{Z'}^2)^2 + M_{Z'}^4 \Gamma_{Z'}^2}
\]

(6.9)
This re-weighting method is validated using dedicated \(Z'\) signal samples generated at different values of \(M_{Z'}\) with un-modified PYTHIA code.

For the analysis at \(\sqrt{s} = 8\) TeV, another technique is used. Instead of re-weighting a dedicated sample of signal events, \(Z'\) signal templates are obtained by re-weighting background \(Z/\gamma^*\) events generated at LO, to match every signal hypothesis considered. The aim is then to replace the matrix elements entering Equation 2.60, the LO differential cross section for \(Z/\gamma^*\) production, by the ones corresponding to \(Z'\) production:

\[
\frac{d\sigma}{dm^2 dy}(pp \rightarrow Z/\gamma^* \rightarrow \mu^+\mu^-) \rightarrow \frac{d\sigma}{dm^2 dy}(pp \rightarrow Z' \rightarrow \mu^+\mu^-) \tag{6.10}
\]

In practice, this is achieved by applying an event-by-event weight to \(Z/\gamma^*\) events as follows:

\[
W(m, q) = \frac{\sum_{i,j \in \{L,R\}} |A_{ij}(Z')|^2}{\sum_{i,j \in \{L,R\}} |A_{ij}(Z/\gamma^*)|^2} \tag{6.11}
\]

where in addition to depending on the invariant mass of the event, the event weight depends on the incoming quarks’ flavour \(q\) via their couplings to the \(Z'\) boson appearing in \(A_{ij}(Z')\). The latter is given by

\[
A_{ij}(Z') = g_{Z'}^2 \frac{g_{\mu j} g'_{\mu j}}{m^2 - M_{Z'}^2 - iM_{Z'} \Gamma_{Z'}} \tag{6.12}
\]

where \(g_{Z'}\) is the universal part of the \(Z'\) coupling and the charges \(g'\) depend on the quantum numbers of the fermions, in accordance with the conventions set in Equation 2.28. The \(Z/\gamma^*\) matrix element \(A_{ij}(Z/\gamma^*)\) is given in Equation 2.56 with \(\hat{s} = m^2\). Example \(Z'_{\text{SSM}}\) signal templates generated with this method are shown in Figure 6.14.

**\(Z'\) signal templates including interference**

The previous strategy can be improved to generate combined templates of \(Z/\gamma^*\) background and \(Z'\) signal, with the interference between the two processes fully taken into account. This implies to substitute the matrix element in the numerator of the event weight as follows:

\[
W'(m, q) = \frac{\sum_{i,j \in \{L,R\}} |A_{ij}(Z/\gamma^*) + A_{ij}(Z')|^2}{\sum_{i,j \in \{L,R\}} |A_{ij}(Z/\gamma^*)|^2} \tag{6.13}
\]

with the \(A_{ij}\) as above. For this analysis, this method is used in the context of Minimal \(Z'\) Models, for which

\[
g_{Z'} g' = g_Y Y + g_{B-L}(B-L) \tag{6.14}
\]
In contrast with the nominal interpretation of the search, where the signal cross section is chosen as the parameter of interest, in the case of Minimal $Z'$ Models the parameters of interest are $\gamma'$ and $\theta_{\text{Min}}$, for this interpretation to remain as general as possible. Therefore, here it is not sufficient to generate signal templates as a function of $m_{\mu^+\mu^-}$ for each $M_{Z'}$ of interest. Two-dimensional templates are necessary: for each value on $M_{Z'}$ and $\theta_{\text{Min}}$ under consideration, combined templates of background and signal are generated as a function of both $m_{\mu^+\mu^-}$ and $\gamma'$. An example of these 2D combined templates is shown in Figure 6.15. In addition to correctly taking into account the interference between $Z/\gamma^*$ and $Z'$ bosons, this approach varies the width $\Gamma_{Z'}$ simultaneously with the signal cross section, as required when varying the couplings between the $Z'$ boson and fermions.

or equivalently, using Equations 3.5 and 3.6:

$$g_{Z'} = \gamma' g_Z$$  \hspace{1cm} (6.15)

$$g' = Y \sin \theta_{\text{Min}} + (B - L) \cos \theta_{\text{Min}}$$  \hspace{1cm} (6.16)

Figure 6.14: Example $Z'_\text{SSM}$ signal templates generated at $\sqrt{s} = 8$ TeV, for $M_{Z'} = 1.5$, 2.5 and 3.5 TeV. The bin width is constant in log $m_{\mu^+\mu^-}$.
Figure 6.15: Example 2D combined template of $Z/\gamma^*$ background and $Z_{\text{Min}}^*$ signal generated at $\sqrt{s} = 8$ TeV, as a function of $\gamma'^4$ and $m_{\mu^+\mu^-}$, for $M_{Z_{\text{Min}}} = 2.5$ TeV and $\theta_{\text{Min}} = 0$.

6.3 Comparison of Data with Background Expectations

This section compares the sum of background expectations to data collected by the ATLAS experiment in 2011 at $\sqrt{s} = 7$ TeV, and in 2012 at $\sqrt{s} = 8$ TeV.

In practice, the search region is initially blinded for $m_{\mu^+\mu^-} > 400$ GeV, in order to prevent the introduction of potential biases in the event selection. The kinematic distributions of individual muon candidates and the dimuon system (with the exception of $m_{\mu^+\mu^-}$) are compared between data and simulation to verify the quality of the simulation on which background estimates rely. The missing transverse energy spectrum is also verified: if large and frequent mis-measurements of muon momentum were present in data and not in the simulation, the $E_T^{\text{miss}}$ spectrum in data would exhibit a shift at higher values. Finally, two-dimensional histograms of $E_T^{\text{miss}}$ vs. leading muon $p_T$ as well as $E_T^{\text{miss}}$ vs. $m_{\mu^+\mu^-}$ are monitored in data, since correlations between these observables are common in events with mis-measured muons.

Only after these verifications, when discrepancies are understood and the frozen selection is approved in the first step of internal review by the collaboration, is the search region fully unblinded. The unblinded distributions are shown here.
6.3.1 Kinematics of the Dimuon System

The dimuon invariant mass spectrum is shown for $m_{\mu^+\mu^-} > 70$ GeV with data at $\sqrt{s} = 7$ TeV in Figure 6.16, and for $m_{\mu^+\mu^-} > 80$ GeV with data at $\sqrt{s} = 8$ TeV in Figure 6.17. Data are compared to the stacked sum of expected backgrounds. Potential contributions from $Z_{SSM}$ bosons are also shown above the expected backgrounds. These histograms are the most important in this analysis, because $m_{\mu^+\mu^-}$ is the discriminating variable used to search for new resonances.

The number of observed events and the corresponding background expected from each source in the search region of the analysis are shown in Tables 6.3 and 6.4 in bins of $m_{\mu^+\mu^-}$. No significant difference between the data and expectations from the Standard Model is observed from these tables; a thorough statistical analysis of the level of agreement with the null hypothesis is carried out in Section 8.1.

Figures 6.18 and 6.19 show the transverse momentum $p_T$ and rapidity $y$ of the dimuon system. The rapidity distributions in data agree well with the expectations from $Z/\gamma^*$ background, however an excess of events in data is seen for both datasets at high dimuon $p_T$. Indeed, while the agreement between data and the simulation is satisfactory in the bulk of the distributions, not enough simulated events remain in the tails, corresponding to a deficit in the amount of radiated gluons with very-high momentum from the initial state of the collision in simulated events. It was verified using the first 1.21 fb$^{-1}$ of data collected at $\sqrt{s} = 7$ TeV that the dimuon $p_T$ is well-modelled when the $Z/\gamma^*$ process is simulated using the ALPGEN generator, which is known to describe events with high jet multiplicities more accurately. This comparison is shown in Figure 6.20.

Ultimately, the results of the search are completely unaffected by this mis-modelling, because the invariant mass of the observed highly-boosted dimuon events is generally on the Z peak or lower. In other words, the mis-modelling in the tail of the dimuon $p_T$ spectrum does not significantly affect the yield predictions in the search region. For the analysis, the $Z/\gamma^*$ samples generated with PYTHIA and POWHEG are chosen because of their significantly higher statistics at high invariant mass.

6.3.2 Kinematics of Individual Muons

Transverse momentum spectra of individual muon candidates are shown in Figures 6.21 and 6.22, separately for the leading and sub-leading muons of each event, i.e. the muons with the highest and second-highest $p_T$. Excesses are observed in the tails of the leading muon momentum distributions in both datasets: this is completely due to the mis-modelling of dimuon $p_T$ and jet multiplicities discussed in the previous section. Indeed, it was verified explicitly that the events responsible for both effects are the same, and that the tail of the muon $p_T$ spectrum is better described by simulation with the ALPGEN generator. This comparison is shown in Figure 6.23.

The angular distributions in the $\eta$ and $\phi$ coordinates are shown in Figure 6.24 for the dataset at $\sqrt{s} = 7$ TeV, and in Figures 6.25 and 6.26 for the dataset at $\sqrt{s} = 8$ TeV. Here the agreement with the Monte Carlo simulation is excellent.
Figure 6.16: Dimuon invariant mass in the selected events at $\sqrt{s} = 7$ TeV.

Table 6.3: Expected and observed number of events for the analysis at $\sqrt{s} = 7$ TeV. The errors quoted include both statistical uncertainties and the systematic uncertainties discussed in Chapter 7.

<table>
<thead>
<tr>
<th>$m_{\mu^+\mu^-}$ [GeV]</th>
<th>110–200</th>
<th>200–400</th>
<th>400–800</th>
<th>800–1200</th>
<th>1200–3000</th>
</tr>
</thead>
<tbody>
<tr>
<td>$Z/\gamma^*$</td>
<td>21200 ± 1200</td>
<td>2090 ± 230</td>
<td>173 ± 15</td>
<td>7.7 ± 0.8</td>
<td>0.98 ± 0.16</td>
</tr>
<tr>
<td>$t\bar{t}, Wt$</td>
<td>900 ± 100</td>
<td>270 ± 50</td>
<td>18 ± 11</td>
<td>0.32 ± 0.07</td>
<td>0.019 ± 0.007</td>
</tr>
<tr>
<td>$WW, WZ, ZZ$</td>
<td>289 ± 32</td>
<td>97 ± 24</td>
<td>11.8 ± 2.7</td>
<td>0.59 ± 0.26</td>
<td>0.087 ± 0.016</td>
</tr>
<tr>
<td>Total</td>
<td>22400 ± 1200</td>
<td>2460 ± 240</td>
<td>203 ± 19</td>
<td>8.7 ± 0.9</td>
<td>1.09 ± 0.16</td>
</tr>
<tr>
<td>Data</td>
<td>21945</td>
<td>2294</td>
<td>197</td>
<td>10</td>
<td>2</td>
</tr>
</tbody>
</table>
Figure 6.17: Dimuon invariant mass in the selected events at $\sqrt{s} = 8$ TeV.

Table 6.4: Expected and observed number of events for the analysis at $\sqrt{s} = 8$ TeV. The errors quoted include both statistical uncertainties and the systematic uncertainties discussed in Chapter 7.

<table>
<thead>
<tr>
<th>$m_{\mu^+\mu^-}$ [GeV]</th>
<th>110–200</th>
<th>200–400</th>
<th>400–800</th>
<th>800–1200</th>
<th>1200–3000</th>
<th>3000–4500</th>
</tr>
</thead>
<tbody>
<tr>
<td>$Z/\gamma^*$</td>
<td>111000 ± 8000</td>
<td>11000 ± 1000</td>
<td>1000 ± 100</td>
<td>49 ± 5</td>
<td>7.3 ± 1.1</td>
<td>0.034 ± 0.022</td>
</tr>
<tr>
<td>$t\bar{t}, Wt$</td>
<td>7100 ± 600</td>
<td>2300 ± 400</td>
<td>160 ± 80</td>
<td>3.0 ± 1.7</td>
<td>0.17 ± 0.15</td>
<td>&lt; 0.001</td>
</tr>
<tr>
<td>WW, WZ, ZZ</td>
<td>1530 ± 180</td>
<td>520 ± 130</td>
<td>64 ± 16</td>
<td>4.2 ± 2.1</td>
<td>0.69 ± 0.30</td>
<td>0.0024 ± 0.0019</td>
</tr>
<tr>
<td>Total</td>
<td>120000 ± 8000</td>
<td>13700 ± 1100</td>
<td>1180 ± 130</td>
<td>56 ± 6</td>
<td>8.2 ± 1.2</td>
<td>0.036 ± 0.023</td>
</tr>
<tr>
<td>Data</td>
<td>120011</td>
<td>13479</td>
<td>1122</td>
<td>49</td>
<td>8</td>
<td>0</td>
</tr>
</tbody>
</table>
Figure 6.18: Dimuon transverse momentum (top) and rapidity (bottom) in the selected events at $\sqrt{s} = 7$ TeV. The bin width of the $p_T$ histogram is constant in $\sqrt{p_T}$. 
Figure 6.19: Dimuon transverse momentum (top) and rapidity (bottom) in the selected events at $\sqrt{s} = 8$ TeV. The bin width of the $p_T$ histogram is constant in $\sqrt{p_T}$. 

\[ \int L \, dt = 20.5 \text{ fb}^{-1} \]
\[ \sqrt{s} = 8 \text{ TeV} \]
Figure 6.20: Dimuon transverse momentum in the first 1.21 fb\(^{-1}\) of data collected at \(\sqrt{s} = 7\) TeV. In the top histogram, the \(Z/\gamma^*\) prediction is evaluated using the nominal generator PYTHIA, while ALPGEN is used for the \(Z/\gamma^*\) prediction in the bottom histogram. The bin width is constant in \(\sqrt{p_T}\).
Figure 6.21: Transverse momentum of the leading muon (top) and sub-leading muon (bottom) in the selected events at $\sqrt{s} = 7$ TeV. The bin width is constant in $\sqrt{p_T}$. 

[Graph showing distribution of transverse momentum for leading and sub-leading muons.]
Figure 6.22: Transverse momentum of the leading muon (top) and sub-leading muon (bottom) in the selected events at $\sqrt{s} = 8$ TeV. The bin width is constant in $\sqrt{p_T}$. 
Figure 6.23: Muon transverse momentum in the first 1.21 fb$^{-1}$ of data collected at $\sqrt{s} = 7$ TeV. In the top histogram, the $Z/\gamma^*$ prediction is evaluated using the nominal generator PYTHIA, while ALPGEN is used for the $Z/\gamma^*$ prediction in the bottom histogram. The bin width is constant in $\sqrt{p_T}$. 
Figure 6.24: $\eta$ and $\phi$ distributions for the selected muons at $\sqrt{s} = 7$ TeV.
Figure 6.25: $\eta$ distributions for the leading muon (top) and sub-leading muon (bottom) in the selected events at $\sqrt{s} = 8$ TeV.
Figure 6.26: $\phi$ distributions for the leading muon (top) and sub-leading muon (bottom) in the selected events at $\sqrt{s} = 8$ TeV.
6.3.3 Missing Transverse Energy Distributions

The missing transverse energy spectrum at $\sqrt{s} = 7$ TeV is shown in Figure 6.27 and the two-dimensional histograms of $E_T^{\text{miss}}$ vs. $m_{\mu^+\mu^-}$ and $E_T^{\text{miss}}$ vs. leading muon $p_T$ observed in data are shown separately for the primary and secondary dimuon selections in Figure 6.28.

The agreement in the $E_T^{\text{miss}}$ distribution between data and the simulation is satisfactory, in spite of the presence of an excess of observed events for $50 \text{ GeV} < E_T^{\text{miss}} < 150 \text{ GeV}$. Perfect agreement is not expected here, since no data quality assessment from the calorimeters is required for this analysis\footnote{Significantly better agreement is observed in dedicated $E_T^{\text{miss}}$ analyses such as Ref. [22] and [28]}; it is more important that no mis-modelling is observed in the tail of the distribution, as this implies that no large systematic mis-measurements of muon momentum are present in data. No correlation is seen in the two-dimensional histograms, with the possible exception of just one outlying event at high-$m_{\mu^+\mu^-}$ in the secondary dimuon selection. The results of the search are unaffected by the presence of this event.

The corresponding histograms for the analysis at $\sqrt{s} = 8$ TeV are shown in Figures 6.29 and 6.30. Here the $E_T^{\text{miss}}$ spectrum is shown separately for the primary and secondary dimuon selections; in both of them, deficits in the number of observed events are present in the bulk of the distributions, but the agreement between data and the simulation is good in the tails. No correlation is observed in the two-dimensional histograms for this dataset.

Figure 6.27: Missing transverse energy in the selected events at $\sqrt{s} = 7$ TeV.
Figure 6.28: Two-dimensional histograms of $E_{\text{miss}}^\mu$ vs. $m_{\mu^+\mu^-}$ (left) and $E_{\text{miss}}^\mu$ vs. leading muon $p_T$ (right) for events passing the primary (top) and secondary (bottom) dimuon selection in collision data at $\sqrt{s} = 7$ TeV.
Figure 6.29: Missing transverse energy in the selected events at $\sqrt{s} = 8$ TeV, for the primary (top) and secondary (bottom) dimuon selection. The bin width is constant in $E_T^{\text{miss}}$. 
Figure 6.30: Two-dimensional histograms of $E_T^{\text{miss}}$ vs. $m_{\mu^+\mu^-}$ (left) and $E_T^{\text{miss}}$ vs. leading muon $p_T$ (right) for events passing the primary (top) and secondary (bottom) dimuon selection in collision data at $\sqrt{s} = 8$ TeV.
Chapter 7

Systematic Uncertainties

7.1 Overview

In this analysis, normalizing all backgrounds to data in the Z peak region simplifies the treatment of systematic uncertainties, as this makes the background estimate insensitive to all uncertainties independent of $m_{\mu^+\mu^-}$, such as the uncertainty on the integrated luminosity. Only the mass-dependent systematic uncertainties on the background prediction therefore need to be considered.

On the other hand, by following this procedure the uncertainty on the signal expectation coming from mass-independent sources is traded for the systematic uncertainty from the theoretical $Z/\gamma^*$ cross section in the normalization region. This is seen most clearly when considering a rearrangement of Equation 6.1 applied to signal processes:

$$
(\sigma B)_{Z'} = \frac{\mu_{Z'}}{(A\varepsilon)_{Z'}} \cdot \frac{1}{L_{int}} \cdot \frac{\mu_{Z} (A\varepsilon)_{Z} (\sigma B)_{Z}}{\mu_{Z'}}
$$

(7.1)

The uncertainty on $(\sigma B)_{Z}$ is propagated accordingly. Following the conventions set in the Exotics working group of the ATLAS collaboration, no other theoretical uncertainty is applied on the signal prediction in the statistical model, because of the model-dependent nature of such calculations.

The main systematic uncertainties to which this analysis is sensitive are listed in Tables 7.1 and 7.2, respectively for the datasets collected in 2011 and 2012. For the analysis at $\sqrt{s} = 7$ TeV, theoretical uncertainties are quoted as a function of truth $m_{\mu^+\mu^-}$, making them overly conservative. This is improved for the analysis at $\sqrt{s} = 8$ TeV, for which uncertainties are quoted with respect to reconstructed $m_{\mu^+\mu^-}$, thereby taking into account muon momentum resolution effects. As it turns out, a large fraction of events in the tail at high $m_{\mu^+\mu^-}$ comes from events with lower truth $m_{\mu^+\mu^-}$, and the relative effect of theoretical uncertainties on the total background yield is in fact smaller.
Table 7.1: Summary of systematic uncertainties on the expected number of events for the search using data collected in 2011 at $\sqrt{s} = 7$ TeV.

<table>
<thead>
<tr>
<th>Source</th>
<th>$m_{\mu^+\mu^-} = 1$ TeV Signal</th>
<th>$m_{\mu^+\mu^-} = 1$ TeV Background</th>
<th>$m_{\mu^+\mu^-} = 2$ TeV Signal</th>
<th>$m_{\mu^+\mu^-} = 2$ TeV Background</th>
<th>$m_{\mu^+\mu^-} = 3$ TeV Signal</th>
<th>$m_{\mu^+\mu^-} = 3$ TeV Background</th>
</tr>
</thead>
<tbody>
<tr>
<td>Normalization</td>
<td>5% N/A</td>
<td>5% N/A</td>
<td>5% N/A</td>
<td>5% N/A</td>
<td>5% N/A</td>
<td>5% N/A</td>
</tr>
<tr>
<td>PDF, $\alpha_s$, scale</td>
<td>N/A 2%</td>
<td>N/A 4.5%</td>
<td>N/A 20%</td>
<td>N/A 44%</td>
<td>N/A 7%</td>
<td>N/A 44%</td>
</tr>
<tr>
<td>Electroweak corr.</td>
<td>3% 6%</td>
<td>3% 6%</td>
<td>3% 6%</td>
<td>3% 6%</td>
<td>3% 6%</td>
<td>3% 6%</td>
</tr>
<tr>
<td>Efficiency</td>
<td>3%</td>
<td>&lt;3%</td>
<td>&lt;3%</td>
<td>&lt;3%</td>
<td>&lt;3%</td>
<td>&lt;3%</td>
</tr>
<tr>
<td>Resolution</td>
<td>3%</td>
<td>&lt;3%</td>
<td>1%</td>
<td>3%</td>
<td>&lt;3%</td>
<td>8%</td>
</tr>
<tr>
<td>Total</td>
<td>6% 8%</td>
<td>8% 21%</td>
<td>10% 46%</td>
<td>10% 46%</td>
<td>10% 46%</td>
<td>10% 46%</td>
</tr>
</tbody>
</table>

Table 7.2: Summary of systematic uncertainties on the expected numbers of events for the search using data collected in 2012 at $\sqrt{s} = 8$ TeV.

<table>
<thead>
<tr>
<th>Source</th>
<th>$m_{\mu^+\mu^-} = 1$ TeV Signal</th>
<th>$m_{\mu^+\mu^-} = 1$ TeV Background</th>
<th>$m_{\mu^+\mu^-} = 2$ TeV Signal</th>
<th>$m_{\mu^+\mu^-} = 2$ TeV Background</th>
<th>$m_{\mu^+\mu^-} = 3$ TeV Signal</th>
<th>$m_{\mu^+\mu^-} = 3$ TeV Background</th>
</tr>
</thead>
<tbody>
<tr>
<td>Normalization</td>
<td>4% N/A</td>
<td>4% N/A</td>
<td>4% N/A</td>
<td>4% N/A</td>
<td>4% N/A</td>
<td>4% N/A</td>
</tr>
<tr>
<td>PDF variation</td>
<td>N/A 5%</td>
<td>N/A 12%</td>
<td>N/A 6%</td>
<td>N/A 12%</td>
<td>N/A 17%</td>
<td>N/A 17%</td>
</tr>
<tr>
<td>PDF choice</td>
<td>N/A &lt;1%</td>
<td>N/A 3%</td>
<td>N/A 6%</td>
<td>N/A 12%</td>
<td>N/A 4%</td>
<td>N/A 4%</td>
</tr>
<tr>
<td>$\alpha_s$</td>
<td>N/A 1%</td>
<td>N/A 3%</td>
<td>N/A 6%</td>
<td>N/A 12%</td>
<td>N/A 4%</td>
<td>N/A 4%</td>
</tr>
<tr>
<td>Electroweak corr.</td>
<td>N/A 1%</td>
<td>N/A 3%</td>
<td>N/A 6%</td>
<td>N/A 12%</td>
<td>N/A 4%</td>
<td>N/A 4%</td>
</tr>
<tr>
<td>$\gamma$-induced corr.</td>
<td>N/A 2%</td>
<td>N/A 3%</td>
<td>N/A 6%</td>
<td>N/A 12%</td>
<td>N/A 4%</td>
<td>N/A 4%</td>
</tr>
<tr>
<td>Beam energy</td>
<td>N/A 2%</td>
<td>N/A 3%</td>
<td>N/A 6%</td>
<td>N/A 12%</td>
<td>N/A 4%</td>
<td>N/A 4%</td>
</tr>
<tr>
<td>Resolution</td>
<td>N/A 2%</td>
<td>N/A 3%</td>
<td>N/A 6%</td>
<td>N/A 12%</td>
<td>N/A 4%</td>
<td>N/A 4%</td>
</tr>
<tr>
<td>Total</td>
<td>4% 7%</td>
<td>4% 15%</td>
<td>4% 23%</td>
<td>4% 23%</td>
<td>4% 23%</td>
<td>4% 23%</td>
</tr>
</tbody>
</table>

than implied by studies carried out at generator-level. Uncertainties having an impact smaller than 3% on the expected number of events are neglected in the statistical model.

7.1.1 Theoretical Uncertainties

The dominant uncertainties in the analysis are theoretical in nature, and come from Parton Distribution Functions (PDFs), introduced in Section 2.2.2. Indeed, $Z/\gamma^*$ production at invariant masses comparable to the beam energy requires both a quark and an antiquark with high momentum fraction $x$. As illustrated in Figure 7.1, because the PDFs of antiquarks are not well-known at such high $x$, the fractional uncertainty on the quark-antiquark luminosity gets large at high $\sqrt{s}/\hat{s} = \sqrt{x_1 x_2}$, resulting in large uncertainties on the background and signal cross sections. Uncertainties due to $\alpha_s$ and QCD scale variations are also considered. The methods used to evaluate these theoretical uncertainties are explained in detail in Sections 7.2 and 7.3.
Figure 7.1: Fractional uncertainty at 90% CL on the quark-antiquark luminosity at √s = 7 TeV due to PDFs, as a function of √s/s = √x1x2. The uncertainty rise at high momentum fraction is clearly visible [167]. Figure credit: G. Watt.

For the analysis on data at √s = 7 TeV, the uncertainty on the electroweak corrections includes uncertainties in the calculation of real boson radiation, as well as the difference in the electroweak scheme definition between PYTHIA and HORACE, in addition to higher-order electroweak and \(\mathcal{O}(\alpha\alpha_s)\) corrections. This treatment was improved for the analysis at √s = 8 TeV: the calculation is performed with higher precision with FEWZ, and uses a consistent electroweak scheme. There a smaller uncertainty is assigned, corresponding to the difference with an independent calculation using the MCSANC [53] generator. The uncertainties on photon-induced corrections are also taken into account in both cases, and are dominated by the uncertainties on the photon PDF and on quark masses.

Systematic uncertainties on sub-leading backgrounds, namely \(t\bar{t}\) and diboson production, are also evaluated as discussed in Section 6.2.1. At high invariant masses the dominant contribution comes from the extrapolation uncertainties. These uncertainties are large with respect to the corresponding backgrounds, but since the \(t\bar{t}\) and diboson backgrounds are an order of magnitude smaller than the \(Z/\gamma^*\) contribution, the resulting uncertainty on the total background is negligible.

7.1.2 Experimental Uncertainties

On the experimental side, the uncertainty on the muon trigger and reconstruction efficiency is dominated by the potential impact of catastrophic energy loss due to muon bremsstrahlung in the
calorimeters or in the MS. While muons detected in the ATLAS experiment are typically minimum-ionizing particles, which is the property that allows them to pass through the calorimeters with little energy loss and reach the Muon Spectrometer, muons at very high momentum enter the regime where radiative energy losses become important, as illustrated in Figure 7.2. In particular, radiative energy losses become larger than energy losses due to ionization for muon energy values above the critical energy $E_{\mu c}$. As shown in Figure 7.3, $E_{\mu c} \sim 150 - 300$ GeV for the iron, copper, tungsten and lead used as absorbing materials in the ATLAS calorimeters, where most of the energy losses occur.

As indicated in Section 5.1, energy depositions are taken into account when extrapolating the tracks measured in the MS to the primary vertex of the event. Nevertheless, if this procedure was not correctly calibrated for muons with very large energy losses, significant differences would be observed between the track parameters measured in the ID and the standalone muon track, used in track matching algorithms. This would in turn negatively impact the muon trigger and reconstruction efficiency, especially for muons at very high momentum for which catastrophic energy losses are more likely.

For the analysis at $\sqrt{s} = 7$ TeV, this efficiency loss is estimated directly from simulation, by quantifying the efficiency as a function of the total energy loss from muons in simulated signal events. Propagating this as a function of $Z'$ mass, the systematic uncertainty is found to be 3%/ TeV. This estimate is improved using more simulated statistics at $\sqrt{s} = 8$ TeV. Looking at muons passing the full selection, no significant decrease in efficiency is observed for energy loss values up to 1 TeV. The systematic uncertainty due to this effect is therefore neglected for this dataset.

The uncertainty on the muon resolution parameters described in Section 5.2 propagates to changes in the background shapes and in the width of signal templates. To evaluate these effects, signal and background estimates are repeated with the muon resolution parameters worsened by their respective uncertainties. The relative differences between the two estimates are then taken as the systematic uncertainties. Figure 7.4 shows the uncertainty on the background estimate due to muon momentum resolution uncertainties for the analysis at $\sqrt{s} = 8$ TeV, while Figure 7.5 demonstrates that signal event migrations due to this effect are negligible. Results are similar at $\sqrt{s} = 7$ TeV.

As discussed in Section 5.1, the muon momentum scale is calibrated using the $Z \rightarrow \mu^+ \mu^-$ peak to a precision of 0.1%. The effect of pileup on the signal acceptance is also negligible: it is checked by varying the number of interactions per proton bunch crossing in simulated signal samples and verifying that the signal acceptance does not change significantly.

Finally, for data collected at $\sqrt{s} = 8$ TeV, a systematic uncertainty of 0.65% on the beam energy of 4 TeV propagates to an uncertainty of 3% on the $Z/\gamma^*$ background yield at high dimuon invariant mass. This effect was not considered for the analysis at $\sqrt{s} = 7$ TeV. The effect on the signal normalization is under 1% for all $Z'$ masses.
Figure 7.2: Stopping power for positive muons in copper as a function of $\beta \gamma = p/Mc$ [52].

Figure 7.3: Muon critical energy for the chemical elements [52].
**Figure 7.4:** Uncertainty on the background estimate due to the muon resolution as a function of $m_{\mu^+\mu^-}$, for the primary dimuon selection (left) and the secondary dimuon selection (right) at $\sqrt{s} = 8$ TeV.

**Figure 7.5:** $Z'$ signal templates at 2 and 3 TeV for the primary dimuon selection at $\sqrt{s} = 8$ TeV, with nominal smearing (blue) and over-smearing increasing the MS resolution smearing constants by their uncertainty (red). The vertical lines indicate the $\pm 1$ RMS ranges corresponding to the nominal templates.

### 7.2 Parton Distribution Function and QCD Uncertainties on Signal Cross Sections

The variation of PDFs and $\alpha_S$ are expected to have a large impact on the $Z'$ cross section as a function of mass. While by convention theoretical uncertainties on signal do not enter the statistical model used in searches for new physics, it is important to be aware of them when converting limits on the signal cross section into model-dependent mass limits. This section explains the techniques used to calculate uncertainties on signal cross sections due to PDF and $\alpha_S$ variations; variants of the same techniques will be applied when quantifying the impact on background estimates in Section 7.3.

In addition to central values of the PDFs, each PDF fit collaboration releases a set of PDFs
varied by their systematic uncertainties. These uncertainties are parametrized using mutually independent parameters, which are called the “eigenvectors” of the PDF set in function space, as they can be varied in orthogonal directions to propagate the systematic uncertainties associated with PDF variations to calculated quantities.

For each PDF eigenvector, the $Z'$ cross section is calculated as a function of mass by generating 100,000 simulated events in PYTHIA. This allows to calculate an asymmetric uncertainty at each mass point using the following equations:

$$\Delta \sigma^+ = \sqrt{\sum_{i=1}^{n} (\max(\sigma^+_i - \sigma_0, \sigma^-_i - \sigma_0))^2}$$  \hspace{1cm} (7.2)

$$\Delta \sigma^- = \sqrt{\sum_{i=1}^{n} (\max(\sigma_0 - \sigma^+_i, \sigma_0 - \sigma^-_i))^2}$$  \hspace{1cm} (7.3)

where $n$ is the number of PDF eigenvectors, $\sigma^+_i$ is the cross section for the higher value of the $i$th PDF eigenvector, $\sigma^-_i$ is the cross section for the lower value of the $i$th PDF eigenvector, and $\sigma_0$ is the cross section for the central value PDF. The larger of the positive and negative variation is then taken as the systematic uncertainty on the $Z'$ cross section.

To generate the MC $Z/\gamma^*$ events used to obtain the $Z'$ signal templates for the analysis using data collected at $\sqrt{s} = 7$ TeV, ATLAS makes use of the modified LO PDF set MRST2007LO**, as described in Section 6.2.1. In contrast with regular LO PDF sets, MRST2007LO** does not include eigenvector variations. The closest LO PDF set, MSTW2008LO, is therefore used to estimate the PDF uncertainties on signal cross sections. In addition to central values, MSTW2008LO has 20 orthogonal eigenvector variations, with high and low values for each eigenvector. The resulting uncertainties are shown in Table 7.3. It is also verified that the $Z'$ cross sections central values calculated using the CTEQ6L1 PDF are within these uncertainties.

The same method is used for the analysis at $\sqrt{s} = 8$ TeV. There, MSTW2008LO is used to obtain the $Z'$ signal templates and for calculating the uncertainties due to PDF. Uncertainties due to $\alpha_S$ variations are also taken into account by calculating cross section values for $\alpha_S$ between 0.11365 and 0.12044, which correspond to the 90% CL $\alpha_S$ limits of MSTW, and taking the extremal variations summed in quadrature with the PDF uncertainty. The combined uncertainties due to PDF and $\alpha_S$ variations are shown in Table 7.4. Here too, it is verified that the $Z'$ cross sections central values calculated using a set from another PDF collaboration, in this case CT10, are within these uncertainties.
Table 7.3: Uncertainty on $Z'$ cross sections due to PDF variations at 90% CL at $\sqrt{s} = 7$ TeV.

<table>
<thead>
<tr>
<th>$Z'$ mass [GeV]</th>
<th>Uncertainty using MSTW2008LO at 90% CL</th>
</tr>
</thead>
<tbody>
<tr>
<td>100</td>
<td>$+3.0%$ $-2.1%$</td>
</tr>
<tr>
<td>200</td>
<td>$+2.6%$ $-2.6%$</td>
</tr>
<tr>
<td>500</td>
<td>$+4.4%$ $-3.7%$</td>
</tr>
<tr>
<td>1000</td>
<td>$+5.5%$ $-7.1%$</td>
</tr>
<tr>
<td>1500</td>
<td>$+8.0%$ $-9.8%$</td>
</tr>
<tr>
<td>2000</td>
<td>$+8.5%$ $-13.3%$</td>
</tr>
</tbody>
</table>

Table 7.4: Uncertainty on $Z'$ cross sections due to PDF and $\alpha_S$ variations at 90% CL at $\sqrt{s} = 8$ TeV.

<table>
<thead>
<tr>
<th>$Z'$ mass [GeV]</th>
<th>Uncertainty using MSTW2008LO at 90% CL</th>
</tr>
</thead>
<tbody>
<tr>
<td>200</td>
<td>$+5.6%$ $-4.7%$</td>
</tr>
<tr>
<td>500</td>
<td>$+4.0%$ $-5.0%$</td>
</tr>
<tr>
<td>1000</td>
<td>$+6.8%$ $-6.7%$</td>
</tr>
<tr>
<td>1500</td>
<td>$+11.0%$ $-10.6%$</td>
</tr>
<tr>
<td>2000</td>
<td>$+17.6%$ $-18.3%$</td>
</tr>
<tr>
<td>2500</td>
<td>$+30.1%$ $-29.7%$</td>
</tr>
<tr>
<td>3000</td>
<td>$+42.5%$ $-42.3%$</td>
</tr>
<tr>
<td>3500</td>
<td>$+51.6%$ $-52.8%$</td>
</tr>
<tr>
<td>4000</td>
<td>$+62.1%$ $-60.5%$</td>
</tr>
<tr>
<td>4500</td>
<td>$+71.1%$ $-71.9%$</td>
</tr>
</tbody>
</table>

7.3 Parton Distribution Function and QCD Uncertainties on the $Z/\gamma^*$ Cross Section

In addition to having an effect on signal cross sections, as described in Section 7.2, PDF variations have an effect on the differential $Z/\gamma^*$ cross section as a function of the dimuon invariant mass $m_{\mu^+\mu^-}$. Indeed this effect represents the dominant source of uncertainty in this analysis.

This section describes in detail how this uncertainty is evaluated, in addition to uncertainties due to QCD scale and $\alpha_S$ variations.
7.3.1 Parton Distribution Function Variations

In a manner identical to the central estimate described in Section 6.2.1, the $Z/\gamma^*$ cross section is calculated at NNLO as a function of $m_{\mu^+\mu^-}$, using the program PHOZPR [114] for the analysis at $\sqrt{s} = 7$ TeV and VRAP [6] for the analysis at $\sqrt{s} = 8$ TeV, for each PDF eigenvector variation, at 90% CL in the MSTW2008NNLO parametrization [132]. The relative deviation of these cross sections from the values calculated using the nominal PDF is interpreted as the uncertainty on the cross section due to each PDF variation. The asymmetric uncertainties thus calculated are shown for each PDF eigenvector in Figures 7.6, 7.7 and 7.8, for the analysis at $\sqrt{s} = 8$ TeV.

Then, it is possible to calculate the total asymmetric uncertainty at each invariant mass point using Equations 7.2 and 7.3. This is the procedure followed for the analysis at $\sqrt{s} = 7$ TeV, and then the larger of the positive and negative asymmetric uncertainties is taken as the total systematic uncertainty on the $Z/\gamma^*$ cross section. This symmetric uncertainty enters the likelihood function discussed in Chapter 8 as a single nuisance parameter.

However, it has been observed that using a single nuisance parameter for the uncertainty due to PDFs can lead to an over-constraint. Since the PDF eigenvectors responsible for the uncertainty at low mass are generally different from the ones responsible for the uncertainty at high mass, the uncertainty values in different mass ranges need to be treated as uncorrelated in the likelihood function. While the most appropriate treatment would then be to assign a distinct nuisance parameter to each of the 20 PDF eigenvectors, such a prescription would drastically increase the dimensionality of the fit. As a result, in addition to introducing potential instabilities in the fit, such an approach would be prohibitive in terms of the computing time required by the statistical framework.

Therefore, the following procedure is implemented for the analysis at $\sqrt{s} = 8$ TeV, as a good approximation of the full statistical treatment. PDF eigenvectors are merged into four PDF eigenvector groups, based on the similar shape of their corresponding uncertainty as a function of $m_{\mu^+\mu^-}$. In the following list, a plus sign means that the eigenvector definition is taken as is, while a minus sign means that the definition is inverted such that the downward eigenvector variation is exchanged with the upward one; this is done so that eigenvectors in a given group behave in the same way.

- Group A consists of eigenvectors 2+, 13+, 14-, 17-, 18+ and 20+. It is dominant nowhere, but its contribution is not negligible.
- Group B consists of eigenvectors 3-, 4-, 9+ and 11+. It is dominant for $m_{\mu^+\mu^-} < 400$ GeV.
- Group C consists of eigenvectors 1+, 5+, 7+, and 8-. It is dominant in the range $400$ GeV < $m_{\mu^+\mu^-} < 1500$ GeV.
- Group D consists of eigenvectors 10+, 12+, 15-, 16- and 19+. It is dominant for $m_{\mu^+\mu^-} > 1500$ GeV.
Figure 7.6: Asymmetric uncertainty on the $Z/\gamma^*$ cross section at $\sqrt{s} = 8$ TeV as a function of $m_{\mu^+\mu^-}$ due to each PDF eigenvector taken separately. Here eigenvectors 1 to 8 are shown.
Figure 7.7: Asymmetric uncertainty on the $Z/\gamma^*$ cross section at $\sqrt{s} = 8$ TeV as a function of $m_{\mu^+\mu^-}$ due to each PDF eigenvector taken separately. Here eigenvectors 9 to 16 are shown.
Figure 7.8: Asymmetric uncertainty on the $Z/\gamma^*$ cross section at $\sqrt{s} = 8$ TeV as a function of $m_{\mu^+\mu^-}$ due to each PDF eigenvector taken separately. Here eigenvectors 17 to 20 are shown.

The remaining eigenvector, number 6, has a shape that does not match any group, and its uncertainty contribution is negligible.

Within a group, the uncertainties from the constituent eigenvectors are combined in the following manner at each mass point:

$$\Delta \sigma^+_G = \text{sign}_G \sqrt{\text{sign}_G \sum_{i=1}^{n_G} \text{sign}(\sigma^+_i - \sigma_0)(\sigma^+_i - \sigma_0)^2}$$  \hspace{1cm} (7.4)

$$\Delta \sigma^-_G = \text{sign}_G \sqrt{\text{sign}_G \sum_{i=1}^{n_G} \text{sign}(\sigma^-_i - \sigma_0)(\sigma^-_i - \sigma_0)^2}$$  \hspace{1cm} (7.5)

where the sum is over the $n_G$ PDF eigenvectors in a given group G, $\sigma^+_i$ is the cross section for the upward variation of the $i^{th}$ PDF eigenvector (downward variation, if inverted), $\sigma^-_i$ is the cross section for the downward variation of the $i^{th}$ PDF eigenvector (upward variation, if inverted), $\sigma_0$ is the cross section for the central value PDF, and sign$_G$ is the sign of the sum inside the square root.

This asymmetric uncertainty on the $Z/\gamma^*$ cross section as a function of $m_{\mu^+\mu^-}$ due to the four
distinct PDF eigenvector groups is shown in Figure 7.9 along with the total symmetric uncertainty obtained using the MSTW prescription (which includes eigenvector 6).

As a closure test, Figure 7.10 shows this same total symmetric uncertainty, compared to the symmetric uncertainty resulting from the addition in quadrature of the uncertainties from the four PDF eigenvector groups. The total symmetric uncertainty as calculated with FEWZ is also shown for comparison. Above the invariant mass cut at 80 GeV used in the analysis at $\sqrt{s} = 8$ TeV, the approximation results in a small overestimate of the total uncertainty, of under 0.35% below 3.5 TeV, and under 1% below 4.5 TeV, comparing calculations done with VRAP. The overestimate is slightly worse when eigenvector 6 is taken into account as a fifth group. The uncertainties obtained with FEWZ and VRAP are in good agreement with each other.

In conclusion, for the analysis on data collected at $\sqrt{s} = 8$ TeV in 2012, the uncertainty due to PDF variations on the $Z/\gamma^*$ cross section as a function of $m_{\mu^+\mu^-}$ is parametrized using four nuisance parameters, corresponding to groups of PDF eigenvectors.

### 7.3.2 Parton Distribution Function Set Choice, QCD Scale and $\alpha_S$

In addition to the uncertainty from PDF variations, we consider uncertainties coming from QCD scale and $\alpha_S$ variations. These are also calculated using PHOZPR for the analysis at $\sqrt{s} = 7$ TeV and VRAP for the analysis at $\sqrt{s} = 8$ TeV. The QCD scale uncertainties are estimated by varying the renormalization and factorization scales simultaneously up and down by a factor of two. The resulting maximum variations are taken as the uncertainties. To estimate the uncertainties due to $\alpha_S$ variations, cross section values as a function of $m_{\mu^+\mu^-}$ are calculated for $\alpha_S$ between 0.11365 and 0.12044, corresponding to the limits on $\alpha_S$ at 90% CL from MSTW2008 [133]. The extremal variations, corresponding to these boundary values of $\alpha_S$, are taken as the asymmetric uncertainty.

For the analysis at $\sqrt{s} = 7$ TeV, the central values for the $Z/\gamma^*$ cross section as a function of $m_{\mu^+\mu^-}$ calculated with different PDF sets are compatible within uncertainties. When repeating this verification at $\sqrt{s} = 8$ TeV by comparing MSTW2008 with the most recent predictions at NNLO and $\alpha_S = 0.117$ from the other PDF fit collaborations CT10 [103, 121], NNPDF2.3 [41], ABM11 [1] and HERAPDF1.5 [113], the central values from ABM11 are found to fall outside of the MSTW2008 PDF uncertainty at 90% CL. An additional uncertainty is therefore assigned to reflect potential differences in the underlying theoretical framework between the PDF fit collaborations. To avoid double-counting the uncertainty due to PDF variations, this additional uncertainty is taken as the difference in quadrature between the cross section central values from ABM11 and the PDF variation uncertainty envelope from MSTW2008.

A summary of the uncertainties on the $Z/\gamma^*$ background estimate is shown in Figure 7.11.
Figure 7.9: Asymmetric uncertainty on the $Z/\gamma^*$ cross section at $\sqrt{s} = 8$ TeV as a function of the dilepton invariant mass due to the four distinct PDF eigenvector groups described in the text. The total symmetric uncertainty is shown in black. The bottom graph zooms in to show the low invariant mass region better.
Figure 7.10: Symmetric uncertainty on the $Z/\gamma^*$ cross section at $\sqrt{s} = 8$ TeV as a function of the dilepton invariant mass resulting from the addition in quadrature of the uncertainties from the four PDF eigenvector groups (in red), compared to the total symmetric uncertainty obtained using the MSTW prescription calculated with VRAP (in black) and FEWZ (in green). The bottom graph zooms in the mass region from 80 to 1500 GeV.
Figure 7.11: Symmetric uncertainty on the $Z/\gamma^*$ cross section at $\sqrt{s} = 8$ TeV as a function of the dilepton invariant mass obtained using the MSTW prescription calculated with VRAP, shown along with the uncertainties due to the QCD scale, $\alpha_S$ variations, photon-induced corrections and higher-order electroweak corrections, as well as the difference between the cross section central values from ABM11 and the PDF variation uncertainty envelope from MSTW, taken as an additional systematic uncertainty. The bottom graph zooms in to show the low invariant mass region better.
Chapter 8

Statistical Methods and Results

This chapter presents the statistical methods used in the search for new neutral high-mass resonances decaying into muon pairs, as well as the results of the analysis. First, the dimuon invariant mass spectra of Figures 6.16 and 6.17 are scanned to look for positive deviations of the observed event yields with respect to the expected backgrounds: significant excesses could be indicative of the existence of a signal. The methods used to quantify this significance are detailed in Section 8.1.

Following the signal search, in the case of a null result, limits are placed on the existence of new phenomena indicative of physics beyond the Standard Model. As indicated in Chapter 3, this search is sensitive to a wide variety of hypotheses, and the interpretations considered here in Sections 8.2 and 8.3 concern models predicting new gauge bosons. These results and many other interpretations were published in Ref. [25] and [36]. Importantly, this search was also performed in the dielectron channel, and published jointly. The results in the $Z'$ interpretation for the combination of both channels are also presented here.

8.1 Signal Search

Two techniques are used to assess the significance of observed excesses in data above background expectations. In both cases this significance is quantified using a $p$-value, defined as the probability, assuming that a signal is absent, of observing an outcome at least as consistent with the existence of signal as the one observed in data. This probability is often quoted in terms of the one-sided integral of a unit-width Gaussian distribution: for example, the integral beyond $+3\sigma$ is equal to $1.35 \times 10^{-3}$. According to the convention used in experimental particle physics, a $p$-value smaller than this number constitutes evidence for the existence of a signal; such evidence is then called a “3σ effect”. Discoveries are only formally claimed when $p < 2.87 \times 10^{-7}$, corresponding to a “5σ effect”.

---

1 In this chapter, the symbol $\ell^\pm$ therefore represents $e^\pm$ and $\mu^\pm$ only, instead of $e^\pm$, $\mu^\pm$ and $\tau^\pm$ as in Chapter 2.

Among other differences the symbol $\mathcal{L}$ is used to represent likelihood functions instead of Lagrangians, etc.
8.1.1 Local $p$-Values

The first method used quantifies the significance individually for each histogram bin of the discriminating variable, in this case the dilepton invariant mass $m_{\ell\ell}$. This local significance is calculated following the methods from Ref. [72].

Using Poisson statistics, the likelihood $\mathcal{L}$ to observe exactly $N$ data events given $\mu$ expected events under hypothesis $H$ is given by

$$\mathcal{L}(N|\mu(H)) = \frac{\mu^N e^{-\mu}}{N!}$$

(8.1)

If $N > \mu$, there is an excess of observed events and the $p$-value then corresponds to the probability to have observed a number $n \geq N$ events given the null hypothesis $H_0$:

$$p(n \geq N|\mu(H_0)) = \sum_{k=N}^{\infty} \frac{\mu^k e^{-\mu}}{k!}$$

$$= 1 - \sum_{k=0}^{N-1} \frac{\mu^k e^{-\mu}}{k!}$$

(8.2)

On the other hand if $N < \mu$, there is an observed deficit with $p$-value

$$p(n \leq N|\mu(H_0)) = \sum_{k=0}^{N} \frac{\mu^k e^{-\mu}}{k!}$$

(8.3)

Systematic uncertainties introduce an overall variance $V$ on the expected number of events $\mu$. This variance is taken into account by setting a gamma distribution as the prior for $\mu$:

$$g(x; a, b) = \frac{b^a}{\Gamma(a)} x^{a-1} e^{-bx}$$

(8.4)

where the parameters $a$ and $b$ are such that

$$\mu = \frac{a}{b} \quad \text{and} \quad V = \frac{a}{b^2} \quad \Rightarrow \quad a = \frac{\mu^2}{V} \quad \text{and} \quad b = \frac{\mu}{V}$$

(8.5)

and $\Gamma$ is the gamma function:

$$\int_0^{\infty} x^{a-1} e^{-bx} dx = \frac{\Gamma(a)}{b^a}$$

(8.6)
Then the likelihood becomes

\[
\mathcal{L}(N|\mu(H), V(H)) = \int_0^{\infty} x^N e^{-x} \frac{1}{N!} g(x; a, b) \, dx
\]

\[
= \int_0^{\infty} x^N e^{-x} \frac{b^a}{N!} \frac{(x)^{a-1} e^{-bx}}{\Gamma(a)} \, dx
\]

\[
= \frac{b^a}{N! \Gamma(a)} \int_0^{\infty} x^{N+a-1} e^{-(1+b)x} \, dx
\]

\[
= \frac{b^a}{N! \Gamma(a)} \Gamma(N+a) (1+b)^{N+a}
\]

and hence the \( p \)-values are

\[
p(n \geq N|\mu(H_0), V(H_0)) = 1 - \sum_{k=0}^{N-1} \mathcal{L}(k|\mu(H), V(H))
\]

\[
p(n \leq N|\mu(H_0), V(H_0)) = \sum_{k=0}^{N} \mathcal{L}(k|\mu(H), V(H))
\]

Figures 8.1 and 8.2 present the dimuon and dielectron invariant mass spectra with data collected in 2011 and 2012 respectively, with the bin-by-bin local significances calculated with this method. The histograms are shown in the search region: \( 128 \, \text{GeV} < m_{\ell\ell} < 3000 \, \text{GeV} \) for the analysis at \( \sqrt{s} = 7 \, \text{TeV} \) and \( 128 \, \text{GeV} < m_{\ell\ell} < 4500 \, \text{GeV} \) for the analysis at \( \sqrt{s} = 8 \, \text{TeV} \). No local excess or deficit with local significance beyond \( \pm 2 \sigma \) is observed in any of the invariant mass histogram bins.

The apparent presence of a global deficit in the dimuon dataset at \( \sqrt{s} = 7 \, \text{TeV} \) was investigated by varying the muon momentum smearing parameters down beyond their uncertainties: this did not significantly impact the observed deficit. This leaves the overall normalization of the background estimates obtained from simulation as a possible cause; this was investigated with the source of this effect remaining unascertained. Such a deficit does not appear in the dataset at \( \sqrt{s} = 8 \, \text{TeV} \).
Figure 8.1: Differences between data and expectation in the and dimuon (top) and dielectron (bottom) channels at $\sqrt{s} = 7$ TeV, with both the statistical and systematic uncertainties taken into account to derive a bin-by-bin local significance, shown in blue.
Figure 8.2: Differences between data and expectation in the dimuon (top) and dielectron (bottom) channels at $\sqrt{s} = 8$ TeV, with both the statistical and systematic uncertainties taken into account to derive a bin-by-bin local significance, shown in blue.
8.1.2 Global \(p\)-Values

While the method discussed in the previous section has the advantage of quantifying potential localized excesses in a general manner, that is independently of any signal hypothesis, another approach is needed to quantify the compatibility of data with specific signal hypotheses, taking the full range of the search region into account. For this purpose, this analysis interfaces the Bayesian Analysis Toolkit (BAT) \[62\] to calculate global \(p\)-values using a template shape fitting technique\(^2\). This technique generalizes the likelihood from Equation 8.1 by taking the product of likelihood functions from all histogram bins of the search region:

\[
\mathcal{L}(N|\mu(H)) = \prod_{j=1}^{n_{\text{bins}}} \frac{\mu_j N_j e^{-\mu_j}}{N_j!}
\]  

(8.10)

where for each bin \(j\) of the search region, \(\mu_j\) represents the expected number of events and \(N_j\) the observed number of events. Such a likelihood function takes into account the full background and signal templates defined in Section 6.2, as well as the complete information from the selected data.

Mass-dependent systematic uncertainties are incorporated into the likelihood function as nuisance parameters \(\theta_s\), which control the magnitude of variations \(\varepsilon_{sj}\) of the expected number of events \(\mu_j\) in bin \(j\) due to each systematic uncertainty \(s\). The nominal range of variations expected from each systematic uncertainty are detailed in Chapter 7. Specifically, the effect of these variations on the expected yields is

\[
\mu_j \rightarrow \tilde{\mu}_j = \mu_j(1 + \sum_{s=1}^{n_{\text{syst}}} \theta_s \varepsilon_{sj})
\]  

(8.11)

The prior on each nuisance parameter is taken to be a Gaussian with zero mean and unit width. The likelihood therefore becomes

\[
\mathcal{L}(N|\mu(H), \theta, \varepsilon) = \prod_{j=1}^{n_{\text{bins}}} \frac{\tilde{\mu}_j N_j e^{-\tilde{\mu}_j}}{N_j!} \prod_{s=1}^{n_{\text{syst}}} e^{-\theta_s^2/2} \frac{1}{\sqrt{2\pi}}
\]  

(8.12)

When multiple search channels are considered, the joint likelihood is the product of the individual likelihoods from each channel. Systematic uncertainties that are fully correlated across channels are assigned the same nuisance parameter, while uncorrelated systematic uncertainties are assigned distinct nuisance parameters in each channel.

To assess the compatibility of experimental data with a specific signal hypothesis \(H_{Z'}\), as compared to the null hypothesis \(H_0\), a test is performed as a function of the \(Z'\) signal cross section \(\sigma_{Z'}\) and mass \(M_{Z'}\), using templates of \(Z'_{SSM}\) signal. The chosen test is based on the Neyman-Pearson lemma \[138\], which applied to this situation states that the most powerful test statistic to reject \(H_0\)

\(^2\) The Bayesian Analysis Toolkit can be used to calculate \(p\)-values, although this concept is intrinsically frequentist.
in favour of $H_{Z'}$ is the Log-Likelihood Ratio (LLR)

$$\text{LLR} = -2 \ln \frac{\mathcal{L}(N|\mu(\hat{\sigma}_{Z'}, \hat{M}_{Z'}), \hat{\theta}, \varepsilon)}{\mathcal{L}(N|\mu(\sigma_{Z'} = 0), \hat{\theta}, \varepsilon)}$$

(8.13)

where $\hat{\sigma}_{Z'}, \hat{M}_{Z'}$ and $\hat{\theta}$ are the best-fit values for the respective parameters given $H_{Z'}$, and $\hat{\theta}$ represents the best-fit values of the nuisance parameters given $H_0$. The best-fit values are the ones at the global maximum of the corresponding likelihood function given the data. Figure 8.3 shows the absolute value of the LLR as a function of the $Z'$ signal mass $M_{Z'}$ and cross section $\sigma_{Z'}$, using the $\sqrt{s} = 7$ TeV dataset, separately for both channels of the search.

The global $p$-value in each channel is then the probability, assuming $H_0$, of observing a value of the LLR test statistic at least as consistent with the existence of signal as the one observed in data, that is

$$p = p(\text{LLR} \leq \text{LLR}_{\text{obs}}|H_0)$$

(8.14)

where because of the minus sign in the definition of the LLR, more negative values imply better agreement with $H_{Z'}$ compared to $H_0$. The LLR distribution given $H_0$ is obtained using pseudo-experiments, i.e. collections of events drawn from the background templates corresponding to the integrated luminosity of the actual dataset. Figure 8.4 shows the LLR distributions obtained using 10,000 pseudo-experiments at $\sqrt{s} = 8$ TeV. To obtain the global $p$-values, these distributions are summed starting from the observed LLR toward more negative values. The $p$-value from each channel is then the value of this sum over the total number of pseudo-experiments considered. The data are found to be consistent with the null hypothesis, with $p$-values of 28% in the dimuon channel and 27% in the dielectron channel. Similar results are obtained in the analysis at $\sqrt{s} = 7$ TeV: there the $p$-values are 68% in the dimuon channel and 36% in the dielectron channel.
Figure 8.3: Absolute value of the LLR used in the search, as a function of the \( Z' \) signal mass \( M_{Z'} \) and cross section \( \sigma_{Z'} \), using the \( \sqrt{s} = 7 \) TeV dataset in the dimuon (top) and dielectron (bottom) channels. The most signal-like value of the LLR found is indicated with a white marker.
Figure 8.4: Distribution of the most signal-like LLR found in each of 10,000 pseudo-experiments at $\sqrt{s} = 8$ TeV, separately for the dimuon (top) and dielectron (bottom) channels. The most signal-like LLR found in data is indicated with the blue arrow, and the $p$-value is the fraction of entries to the left of this arrow.
8.2 Limits on $Z'_{SSM}$ and $E_6$ $Z'$ Bosons

Since the data are found to be consistent with predictions from the Standard Model, limits are set on the existence of $Z'$ bosons, using BAT. This section presents limits set on the cross section $\sigma$ for $Z'$ boson production times its branching fraction $B$ in individual dilepton channels.

As explained in Section 6.2.3, in this nominal approach to the analysis each choice of $Z'$ boson type and mass $M_{Z'}$ corresponds to a specific signal template, and the expected number of events $\mu(H_{Z'})$ only depends on $\sigma B$. The general expression for the binned Poisson likelihood function in Equation (8.12) can then be reduced to a function of the parameter of interest, in this case $\sigma B$, by performing a numerical integration over the nuisance parameters using a Markov chain Monte Carlo algorithm. The marginalized likelihood function is thus

$$L'(\sigma B) = \int L(N|\sigma B, \theta, \varepsilon) \prod_{s=1}^{n_{\text{sys}}} d\theta_s$$

(8.15)

Using this function, the posterior probability density $L'(\sigma B|N)$ is obtained using Bayes’ theorem:

$$L'(\sigma B|N) = L'(N|\sigma B) \frac{\pi(\sigma B)}{\pi(N)}$$

(8.16)

where $\pi(\sigma B)$ is the prior for $\sigma B$, taken as a constant, and $\pi(N)$ is the prior for the observed number of events, which does not depend on $\sigma B$. The most probable signal strength given the data is given by the maximum of $L'(\sigma B|N)$.

Limits at 95% CL are placed on $\sigma B$ by integrating the posterior probability density as follows, to find the limit value $(\sigma B)_{95}$ such that

$$0.95 = \frac{\int_0^{(\sigma B)_{95}} L'(\sigma B|N)d(\sigma B)}{\int_0^{\infty} L'(\sigma B|N)d(\sigma B)}$$

(8.17)

Figure 8.5 shows the observed limits on $\sigma B$ for $Z'_{SSM}$ and $E_6$ $Z'$ production calculated in this manner as a function of $M_{Z'}$, using data collected at $\sqrt{s} = 7$ TeV for the dimuon channel and for the combined dielectron and dimuon channels. The range of expected limits calculated using pseudo-experiments is also shown for comparison. The limits using data collected at $\sqrt{s} = 8$ TeV are shown in Figure 8.6 for $E_6$ $Z'$ production, and Figure 8.7 for $Z'_{SSM}$ production.

Lower limits on the mass of $Z'$ bosons are obtained from the upper limits on $\sigma B$ by finding their intersection with the theory curves for $\sigma B$ as a function of $M_{Z'}$ for the different hypotheses under consideration. Table 8.1 shows the observed and expected values of these mass limits for the combination of dielectron and dimuon channels. The ratio of combined limits on $\sigma(Z'_{SSM} \rightarrow \ell^+ \ell^-)$ to the theoretical cross section is shown in Figure 8.8 for all four published results of this search by the ATLAS collaboration [18, 19, 25, 36].

146
Figure 8.5: Observed (red line) and expected upper limits on $\sigma B$ for $Z'$ boson production at $\sqrt{s} = 7$ TeV for the dimuon channel (top) and for the combined dielectron and dimuon channels (bottom). The median expected limit is shown as a black dashed line, with the ranges at $\pm 1\sigma$ and $\pm 2\sigma$ around the median shown as the yellow and green bands. The region above each of these lines is excluded at 95% CL. The three theory curves for $Z'_{\text{SSM}}$ and $E_6$ $Z'$ boson production are also shown. Other theory curves belonging to $E_6$-motivated models fall between the $Z'_{\psi}$ and $Z'_{\chi}$, since these signals respectively have the smallest and largest $\sigma B$. The dashed lines around the $Z'_{\text{SSM}}$ theory curve represent the theoretical uncertainty, which is similar for the other theory curves.
Table 8.1: Observed and expected lower limits at 95% CL on the mass of $Z'_{\text{SSM}}$ and $E_6 Z'$ bosons obtained by the ATLAS experiment.

<table>
<thead>
<tr>
<th>Centre-of-mass energy</th>
<th>Channel</th>
<th>Mass limits at 95% CL [TeV]</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td></td>
<td>$Z'_{\text{SSM}}$</td>
</tr>
<tr>
<td>$\sqrt{s} = 7 \text{ TeV}$</td>
<td>$pp \to e^+e^-, \mu^+\mu^-$</td>
<td>2.22 (2.25)</td>
</tr>
<tr>
<td>$\sqrt{s} = 8 \text{ TeV}$</td>
<td>$pp \to e^+e^-, \mu^+\mu^-$</td>
<td>2.90 (2.87)</td>
</tr>
</tbody>
</table>

Figure 8.6: Observed (red line) and expected upper limits on $\sigma B$ for the $E_6 Z'_{\psi}$ boson production at $\sqrt{s} = 8 \text{ TeV}$ for the combination of the dielectron and dimuon channels. The median expected limit is shown as a black dashed line, with the ranges at $\pm 1 \sigma$ and $\pm 2 \sigma$ around the median shown as the yellow and green bands. The region above each of these lines is excluded at 95% CL. The two theory curves corresponding to $Z'_{\psi}$ and $Z'_{\chi}$ boson production are also shown. Other $Z'$ signals belonging to $E_6$-motivated Models fall between these two cases, since $Z'_{\psi}$ and $Z'_{\chi}$ respectively have the smallest and largest $\sigma B$. The thickness of the $Z'_{\psi}$ theory curve represents the theoretical uncertainty, which is similar for the $Z'_{\chi}$ theory curve.
Figure 8.7: Observed and median expected upper limits on $\sigma B$ for $Z_{SSM}'$ boson production at $\sqrt{s} = 8$ TeV for the exclusive dimuon and dielectron channels, and for both channels combined. The region above each of these lines is excluded at 95% CL. The $Z_{SSM}'$ theory curve is also shown; the grey area represents its theoretical uncertainty.

Figure 8.8: Ratio of the observed limits for the $Z_{SSM}'$ search to the $Z_{SSM}'$ cross section times branching fraction for the combination of dielectron and dimuon channels. The region above each line is excluded at 95% CL. All four published results of this search by the ATLAS collaboration are shown.
8.3 Limits on Minimal $Z'$ Models

In the context of Minimal $Z'$ Models, limits are placed not on the cross section times branching fraction $\sigma B$, but rather on the $Z'$ coupling $\gamma'$ itself as a function of the mass $M_{Z'}$ and of the mixing angle $\theta_{\text{Min}}$.

As indicated in Section 6.2.3, in this case signal templates are generated as a function of both $m_{\mu^+\mu^-}$ and $\gamma'$ for each pair of values of $M_{Z'}$ and $\theta_{\text{Min}}$ under consideration. The likelihood function is otherwise defined analogously: for each value of $M_{Z'}$ and $\theta_{\text{Min}}$, the expected number of events $\mu_j(H)$ and systematic variations $\epsilon_{sj}$ now depend on both $m_{\mu^+\mu^-}$ and $\gamma'$ instead of only on $m_{\mu^+\mu^-}$. The marginalized likelihood function becomes

$$\mathcal{L}'(N|\gamma') = \int \mathcal{L}(N|\gamma', \theta, \epsilon) \prod_{s=1}^{n_{\text{syst}}} d\theta_s$$

Bayes’ theorem is used with a prior $\pi(\gamma')$ constant in $\gamma'$: this is justified by the fact that the cross section for $Z'$ boson production is proportional to $\gamma'^4$. Limits at 95% CL are then placed on $\gamma'$ by integrating the posterior probability density in the same way as in the previous section, to find $\gamma'_95$ such that

$$0.95 = \frac{\int_0^{\gamma'_95} \mathcal{L}'(\gamma'|N) d\gamma'}{\int_0^{\infty} \mathcal{L}'(\gamma'|N) d\gamma'}$$

The resulting upper limits on $\gamma'$ are shown as a function of $M_{Z'}$ in Figure 8.9 for the dataset at $\sqrt{s} = 7$ TeV and in Figure 8.10 for the dataset at $\sqrt{s} = 8$ TeV. In the latter case, limits are also displayed as a function of $\theta_{\text{Min}}$ for specific values of $M_{Z'}$ in Figure 8.11.

Lower limits on the $Z'$ boson mass can also be obtained in this framework, by finding the intersection of the upper limits on $\gamma'$ with the nominal $\gamma'$ values from theory given in Table 3.1, with $\sin \theta_W = 0.48$. The results are given in Table 8.2. The ranges of observed and expected limits on $M_{Z'}$ for $\theta_{\text{Min}} \in [0, \pi]$ and representative values of $\gamma'$ are given in Table 8.3 and the ranges of observed and expected limits on $\gamma'$ for $\theta_{\text{Min}} \in [0, \pi]$ and representative values of $M_{Z'}$ are given in Table 8.4.
Figure 8.9: Observed and expected limits on $\gamma'$ as a function of the $Z'$ mass at $\sqrt{s} = 7$ TeV, for the dimuon channel (top) and for the combined dielectron and dimuon channels (bottom). Two limit curves are displayed for representative values of $\theta_{\text{Min}}$, which at specific values of $\gamma'$ correspond to the $Z'_{R}$ and $Z'_{B-L}$ models. The collection of all limit curves for $\theta_{\text{Min}} \in [0, \pi]$ forms the grey band. The region of parameter space above each limit curve is excluded at 95% CL.
Figure 8.10: Observed and expected limits on $\gamma'$ as a function of the $Z'$ mass at $\sqrt{s} = 8$ TeV, for the dimuon channel (top) and for the combined dielectron and dimuon channels (bottom). Three limit curves are displayed for representative values of $\theta_{\text{Min}}$, which at specific values of $\gamma'$ correspond to the $Z'_{\chi}$, $Z'_{R}$ and $Z'_{B-L}$ models. The collection of all observed limit curves for $\theta_{\text{Min}} \in [0, \pi]$ forms the grey band, while the grey dotted lines delimit the area containing the corresponding expected limit curves. The region of parameter space above each limit curve is excluded at 95% CL.
Figure 8.11: Observed and expected limits on $\gamma'$ as a function of $\theta_{\text{Min}}$ at $\sqrt{s} = 8$ TeV, for the dimuon channel (top) and for the combined dielectron and dimuon channels (bottom). The limits are set for seven representative values of $M_{Z'_{\text{Min}}}$. The region of parameter space above each limit curve is excluded at 95% CL.
Table 8.2: Observed and expected lower limits at 95% CL on the mass of $Z'_{B-L}$, $Z'_X$ and $Z'_R$ bosons obtained by the ATLAS experiment in the context of Minimal $Z'$ Models. Here interference effects are taken into account and the width $\Gamma_{Z'}$ is varied simultaneously with $\sigma_{Z'}$, resulting in slightly weaker limits for $Z'_X$ than with the nominal approach.

<table>
<thead>
<tr>
<th>Centre-of-mass energy</th>
<th>Channel</th>
<th>Mass limits at 95% CL [TeV]</th>
<th>$Z'_{B-L}$ obs. (exp.)</th>
<th>$Z'_X$ obs. (exp.)</th>
<th>$Z'_R$ obs. (exp.)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\sqrt{s} = 7$ TeV</td>
<td>$pp \to e^+e^-, \mu^+\mu^-$</td>
<td>2.13 (2.14)</td>
<td>1.89 (1.93)</td>
<td>2.11 (2.11)</td>
<td></td>
</tr>
<tr>
<td>$\sqrt{s} = 8$ TeV</td>
<td>$pp \to e^+e^-, \mu^+\mu^-$</td>
<td>2.80 (2.76)</td>
<td>2.51 (2.46)</td>
<td>2.73 (2.69)</td>
<td></td>
</tr>
</tbody>
</table>

Table 8.3: Range of the observed and expected lower limits at 95% CL on the $Z'_{\text{Min}}$ boson mass for $\theta_{\text{Min}} \in [0, \pi]$ and representative values of the relative coupling strength $\gamma'$. Both lepton channels are combined.

<table>
<thead>
<tr>
<th>Centre-of-mass energy</th>
<th>$\gamma'$</th>
<th>Range of limits on $M_{Z'_{\text{Min}}}$ [TeV]</th>
<th>$Z'_{\text{Min}}$ obs. (exp.)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\sqrt{s} = 7$ TeV</td>
<td>0.1</td>
<td>0.67-1.43 (0.58-1.47)</td>
<td></td>
</tr>
<tr>
<td></td>
<td>0.2</td>
<td>1.11-2.10 (1.17-2.07)</td>
<td></td>
</tr>
<tr>
<td>$\sqrt{s} = 8$ TeV</td>
<td>0.2</td>
<td>1.24-2.28 (1.22-2.39)</td>
<td></td>
</tr>
<tr>
<td></td>
<td>0.3</td>
<td>1.89-2.93 (1.83-2.82)</td>
<td></td>
</tr>
<tr>
<td></td>
<td>0.4</td>
<td>2.09-3.12 (2.03-3.08)</td>
<td></td>
</tr>
<tr>
<td></td>
<td>0.5</td>
<td>2.20-3.24 (2.16-3.19)</td>
<td></td>
</tr>
<tr>
<td></td>
<td>0.6</td>
<td>2.31-3.39 (2.24-3.32)</td>
<td></td>
</tr>
</tbody>
</table>

Table 8.4: Range of the observed and expected upper limits at 95% CL on the relative coupling strength $\gamma'$ for $\theta_{\text{Min}} \in [0, \pi]$ and representative values of the $Z'_{\text{Min}}$ boson mass. Both lepton channels are combined.

<table>
<thead>
<tr>
<th>Centre-of-mass energy</th>
<th>$M_{Z'_{\text{Min}}}$ [TeV]</th>
<th>Range of limits on $\gamma'$</th>
<th>$Z'_{\text{Min}}$ obs. (exp.)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\sqrt{s} = 7$ TeV</td>
<td>1.0</td>
<td>0.08-0.16 (0.07-0.15)</td>
<td></td>
</tr>
<tr>
<td></td>
<td>2.0</td>
<td>0.16-1.10 (0.17-1.01)</td>
<td></td>
</tr>
<tr>
<td>$\sqrt{s} = 8$ TeV</td>
<td>1.0</td>
<td>0.08-0.16 (0.08-0.16)</td>
<td></td>
</tr>
<tr>
<td></td>
<td>2.0</td>
<td>0.17-0.34 (0.15-0.37)</td>
<td></td>
</tr>
<tr>
<td></td>
<td>3.0</td>
<td>0.33-1.71 (0.35-1.90)</td>
<td></td>
</tr>
</tbody>
</table>
Looking forward, Run-II of the LHC in 2015–2017 will be a very exciting opportunity to bring this search to new heights. Increasing the centre-of-mass energy from 8 TeV to at least 13 TeV will drastically improve the search prospects at high resonance masses: to give a specific example, the parton luminosity for quark-antiquark processes is expected to grow by a factor \( O(100) \) for a resonance mass of 4 TeV, as shown in Figure 9.1. As a result, the current limits on new neutral resonances are expected to be surpassed with as little as 5 fb\(^{-1}\) of Run-II data. By the end of Run-II, with an integrated luminosity of 75 to 100 fb\(^{-1}\), the reach of the search will be extended to even lower signal cross sections across the whole invariant mass spectrum. Then, a total of up to 300 fb\(^{-1}\) is expected in Run-III of the LHC, and up to 3000 fb\(^{-1}\) will be delivered at the High-Luminosity LHC (HL-LHC) by the year 2030 [26]. This will make possible a direct observation of the \( H \rightarrow \mu^+ \mu^- \) decay channel, thereby allowing to measure the Yukawa coupling of the Higgs boson to muons, in addition to any unexpected discoveries to come. Figure 9.2 shows the expected limits on \( Z' \) boson production in the dimuon channel with the projected HL-LHC dataset: the lower limit at 95% CL on the mass of \( Z'_{\text{SSM}} \) bosons is expected to reach 7.6 TeV.

Such improvements will not come without hard work by the collaboration. The increased event pileup coming from up to 70 interactions per proton-proton bunch crossing on average in Run-III, and up to 140 interactions per crossing on average at the HL-LHC, will bring tremendous gains in instantaneous luminosity but also challenges in triggering and reconstructing these events. In particular, the Inner Detector of ATLAS will need to be completely replaced to accommodate the higher charged track rates and number of collision vertices to identify. The calorimeters and the Muon Spectrometer will be upgraded in the forward region of the detector, in the first case to ensure that the liquid-argon calorimeter performance remains acceptable amid larger radiation levels, and in the second to mitigate the effects of increased fake muon backgrounds on the tracking efficiency and trigger rates. As well, new front-end electronics will be necessary in order to read out the information from all detector subsystems faster, which along with developments in the trigger software...
will considerably improve trigger decision times. Dedicated research and development is ongoing to ensure that the ambitious goals of the LHC machine are met by corresponding upgrades to its detectors, in order to enable their operation for years to come.

Figure 9.1: Ratios of the parton luminosity accessible at the LHC at $\sqrt{s} = 13$ TeV compared to that at $\sqrt{s} = 8$ TeV for gluon-gluon, gluon-quark and quark-antiquark processes. Figure credit: W. J. Stirling.

Figure 9.2: Expected upper limits on $\sigma B$ for $Z'_{SSM}$ boson production for the projected HL-LHC dataset in the dimuon channel. The median expected limit is shown as a black dashed line, with the ranges at $\pm 1\sigma$ and $\pm 2\sigma$ around the median shown as the yellow and green bands. The region above each of these lines is excluded at 95% CL. The $Z'_{SSM}$ theory curve is also shown [26].
Chapter 10

Conclusion

In the search for new neutral high-mass resonances decaying into muon pairs performed by the ATLAS experiment, the dimuon invariant mass spectrum was compared to expectations from the Standard Model of particle physics. A wide variety of hypotheses beyond the Standard Model predict the existence of a resonant peak in this spectrum corresponding to a new particle with a mass larger than the $Z$ boson mass.

A detailed understanding of the performance of the detector for reconstructing muons at very high momentum is absolutely necessary for this search to be successful, since only the best-measured reconstructed muon tracks from the detector have a momentum resolution that makes possible the reconstruction of a peak at the TeV scale. It is also crucial to prevent contamination of the signal region from mis-measured tracks that falsely appear at high momentum. Therefore, stringent selection criteria were defined for the muon tracks used in the search: for example, most of the selected muons are well-measured in all three stations of the Muon Spectrometer of ATLAS. The signal acceptance of the search was increased following dedicated studies, which found muon tracks passing through two stations of the Muon Spectrometer but with acceptable momentum resolution. Other selection criteria on the isolation and impact parameter of muon tracks eliminated the backgrounds due to physical processes with muons from hadronic jets and cosmic rays.

Following the event selection, the observed event yields were compared with state-of-the-art theoretical calculations for the expected dimuon invariant mass spectrum from the Standard Model, with the acceptance, efficiency and resolution of the detector for reconstructing muon tracks taken into account using Monte Carlo simulations. Systematic uncertainties on this prediction were carefully evaluated. The dominant uncertainty comes from our limited knowledge of PDFs at the high momentum fraction values necessary for the production of a resonance at high mass.

The agreement with the predictions from the Standard Model is remarkable. No significant excess of observed events was found in the full proton-proton collision dataset collected at $\sqrt{s} = 7$ TeV and $\sqrt{s} = 8$ TeV in Run-I of the LHC, and limits were consequently placed on the existence of new...
gauge bosons $Z'$. These new gauge bosons are predicted in theoretical hypotheses that reach beyond the Standard Model in order to address some of its shortcomings. In particular, limits were set at 95% CL excluding $Z'$ bosons from the Sequential Standard Model with masses below 2.90 TeV, and $Z'$ from Grand Unified Theories with masses below 2.51 to 2.80 TeV depending on the symmetry breaking scenario. Limits on $Z'$ bosons set by the CMS collaboration are similar. These results represent an excellent improvement over previous limits, set around 1 TeV by the CDF and D0 experiments at the Tevatron.

Furthermore, alternative theoretical hypotheses predict other types of new neutral particles that can also decay into muon pairs, and this search is also sensitive to them. Public results by the ATLAS collaboration [25, 36] place limits on a wide variety of these models using exactly the same methods as detailed in this dissertation. As well, it is possible to recast these experimental results in order to set bounds relevant to other hypotheses. The results of this analysis therefore help to guide future theoretical research by restricting the allowed parameter space for new models.

The sensitivity of this search will keep improving in Run-II of the Large Hadron Collider and beyond, as the energy and intensity frontiers are pushed further by technological developments. The discovery of a new neutral resonance would have a profound impact on our understanding of the universe. Should such a new particle be observed, an analysis of the angular distributions of the final-state leptons would provide key information towards identifying its properties. Otherwise, in the case of a null result, increasing the exclusion limits will provide even more stringent bounds on physics beyond the Standard Model. Either way, the continuation of this search, and more generally of the physics program of the ATLAS and CMS experiments at the LHC, will contribute in fundamental ways to the advancement of knowledge about the elementary constituents of the world we live in.
Bibliography


[25] ATLAS Collaboration. Search for high mass resonances decaying to dilepton final states in

[26] ATLAS Collaboration. Physics at a high-luminosity LHC with ATLAS. 2013,

[27] ATLAS Collaboration. Improved luminosity determination in pp collisions at $\sqrt{s} = 7$ TeV

[28] ATLAS Collaboration. Performance of missing transverse momentum reconstruction in
ATLAS studied in proton-proton collisions recorded in 2012 at $\sqrt{s} = 8$ TeV. Technical

[29] ATLAS Collaboration. Measurements of Higgs boson production and couplings in diboson


[31] ATLAS Collaboration. Evidence for Higgs boson decays to the $\tau^+\tau^-$ final state with the

[32] ATLAS Collaboration. Preliminary results on the muon reconstruction efficiency,
momentum resolution, and momentum scale in ATLAS 2012 pp collision data. Technical
Report ATLAS-CONF-2013-088, CERN, 2013. → pages 64, 69, 70, 71, 72, 73

[33] ATLAS Collaboration. ATLAS Experiment – Public Results, 2014. URL

[34] ATLAS Collaboration. Updated coupling measurements of the Higgs boson with the
ATLAS detector using up to 25 fb$^{-1}$ of proton-proton collision data. Technical Report

[35] ATLAS Collaboration. Muon reconstruction efficiency and momentum resolution of the
ATLAS experiment in proton-proton collisions at $\sqrt{s} = 7$ TeV in 2010. *Submitted to

[hep-ex]. → pages 2, 137, 146, 158


→ pages 28

→ pages 27

→ pages 10

→ pages 6

→ pages 33

→ pages 28

→ pages 28

→ pages 22

→ pages 91, 94

→ pages 201

→ pages 95, 133

→ pages 91

→ pages 29

→ pages 27


169


Appendix A

Search for Vector Boson Fusion

\( H \to WW^* \to \ell\nu\ell\nu \)

This appendix discusses the multivariate analysis looking for Vector Boson Fusion (VBF) production of Higgs bosons decaying into W boson pairs in the \( \ell\nu\ell\nu \) final states. After a brief introduction to the status of Higgs boson observations at the LHC in Section A.1 and a description of the analysis goals and strategy in Section A.2, specific contributions to the upcoming result are presented in Section A.3 and future perspectives are outlined in Section A.4.

A.1 Status of Higgs Boson Observations at the Large Hadron Collider

Following the discovery of a Higgs boson at the LHC in 2012 [23, 76], physicists started measuring its mass, spin, parity and couplings to the other particles. So far, the mass of the new particle has been measured to be \( 125.5 \pm 0.2 \) (stat.) \( ^{+0.5}_{-0.6} \) (syst.) GeV by the ATLAS collaboration [29] and \( 125.6 \pm 0.4 \) (stat.) \( \pm 0.2 \) (syst.) GeV by the CMS collaboration [80]. As well, the data favour the 0\(^{+}\) spin-parity hypothesis over the 0\(^{-}\), 1\(^{+}\), 1\(^{-}\) and 2\(^{+}\) hypotheses at 95% CL or higher [30, 80].

The couplings of the new particle are obtained from measurements of its production and decay rates: it is therefore important to detect it in as many channels as possible. The four production channels with the highest cross sections are, in order, gluon fusion, VBF, associate production with a vector boson, and top quark fusion. Their relative size is shown as a function of the Higgs boson mass \( M_H \) in Figure A.1. On the other side, the cross section times branching fraction to observable final states following a Higgs boson decay are shown in Figure A.2.

Of all the different combinations of initial and final states, only two have been decidedly observed so far: gluon fusion production with decays to \( \gamma\gamma \) [29, 76], and \( ZZ^* \to \ell^+\ell^-\ell^+\ell^- \) [29, 80].

\(^1\) In this appendix, the symbol \( \ell^\pm \) represents \( e^\pm \) or \( \mu^\pm \) only, with \( \tau^\pm \) indicated explicitly.
Figure A.1: Higgs boson production cross section by channel as a function of $M_H$. Figure credit: R. Tanaka.

Figure A.2: Higgs boson cross section times branching fraction to observable final states as a function of $M_H$. Figure credit: R. Tanaka.
There is also very strong evidence for gluon fusion production decaying to $WW^* \rightarrow \ell\nu\ell\nu$ [29, 79] and for VBF production decaying in the $\tau^+\tau^-$ final state [31, 81]. Figure A.3 shows the histograms from the ATLAS collaboration corresponding to these four results; similar histograms are available from the CMS collaboration.

Figure A.3: Histograms with data from the ATLAS experiment displaying evidence for the production of a Higgs boson in four different channels: gluon fusion production followed by decays to $\gamma\gamma$ [29] (top left), $ZZ^* \rightarrow \ell^+\ell^-\ell^+\ell^-$ [29] (top right) and $WW^* \rightarrow \ell\nu\ell\nu$ [29] (bottom left), as well as VBF production decaying in the $\tau^+\tau^-$ final state [31] (bottom right).
A.2 Analysis Goals and Strategy

The goal of this analysis is to search for Vector Boson Fusion $H \rightarrow WW^* \rightarrow \ell \nu \ell \nu$. Measuring the production and decay rate in this channel is particularly interesting, because the Higgs boson couples to $W$ bosons in both the initial and the final states, as illustrated in Figure A.4. Although $Z$ bosons also contribute in the initial state, the production is dominated by $W$ boson fusion. Such a measurement is therefore particularly sensitive to the coupling $\kappa_W$ between the Higgs and $W$ bosons. Explicitly:

$$\sigma(\text{VBF } H \rightarrow WW^*) \propto \kappa_W^4 \Rightarrow \frac{\Delta \sigma}{\sigma} = 4 \frac{\Delta \kappa_W}{\kappa_W}$$

which implies, for example, that achieving a relative uncertainty of 60% on the measurement of the cross section $\sigma$ would constrain the coupling $\kappa_W$ to a relative uncertainty of only 15%.

The decays of $WW^*$ to $\ell \nu \ell \nu$ have small branching fractions, because of the smaller branching fraction of $W$ bosons to leptonic final states than to hadronic final states. They are nevertheless the most promising ones, because the background rates are drastically reduced, especially from $W + \text{jets}$ and QCD multi-jet processes. In particular, the Different-Flavour (DF) channel, where one charged lepton is an electron and the other is a muon, is the principal search channel because background contributions from $Z/\gamma^* \rightarrow \ell^+ \ell^-$ are absent. The channels with two electrons or two muons in the final state are merged, and named the Same-Flavour (SF) channel.

Analyses looking for $H \rightarrow WW^* \rightarrow \ell \nu \ell \nu$ have already been performed by the ATLAS and CMS experiments. Based on the full Run-I dataset, the CMS collaboration has observed a signal strength normalized to the expectation from the Standard Model (SM) of $\sigma/\sigma_{SM} = 0.62^{+0.58}_{-0.47}$, corresponding to a significance of 1.3 standard deviations (2.1 expected) [79]. On the other hand, following a first analysis of the full Run-I dataset, the ATLAS collaboration has observed $\sigma/\sigma_{SM} = 1.4 \pm 0.7$, with a significance of 2.5 standard deviations (1.6 expected) [29]. Histograms showing the data and expected yields corresponding to this last result are shown in Figure A.5. The discriminant variable

Figure A.4: Feynman diagram for VBF $H \rightarrow WW^*$ at Leading Order.
used in the final stage of the search is the transverse mass $m_T$:

$$m_T = \sqrt{(E_T^{\ell\ell} + E_T^{\text{miss}})^2 - |p_T^{\ell\ell} + p_T^{\text{miss}}|^2}, \text{ where } E_T^{\ell\ell} = \sqrt{|p_T^{\ell\ell}|^2 + m_{\ell\ell}^2}$$

(A.2)

Figure A.5: Transverse mass $m_T$ used as discriminant in the final stage of the cut-based search for VBF $H \rightarrow WW^{*} \rightarrow \ell\nu\ell\nu$, for the DF channel (top) and the SF channel (bottom) [29].

The search described here aims to re-analyze the data collected in Run-I by the ATLAS experiment using multivariate analysis techniques. Such techniques are especially appropriate when looking for a small signal among a large variety of background processes, and many discriminat-
ing variables are available. Multivariate analysis exploits the power of correlations between these variables to achieve additional separation between signal and background.

The remainder of this section first describes the physical objects that are observable in the final state of interest in Section A.2.1, and the background processes that yield the same final-state objects in Section A.2.2. Section A.2.3 discusses the unique topological features of the signal process, linked to the variables that make its separation from background possible. Finally, the multivariate analysis techniques investigated are described: Boosted Decision Trees (BDTs) are explained in Section A.2.4, and the Matrix Element (ME) method in Section A.2.5. Of the two, only the method using BDTs is used for the analysis, primarily because the ME method was found to be too computationally demanding.

A.2.1 Physical Objects

In this analysis, the final state of interest is diverse: two light-quark jets accompany the two $W$ bosons, which decay into final states with two electrons, two muons, or one electron and one muon, always with real missing transverse momentum ($E_T^\text{miss}$) caused by the accompanying neutrinos. Because the charged leptons come from $W$ decays, their expected momentum is relatively low and thus easier to measure well. Further, the two jets are expected to be very forward, often requiring the use of the Forward Calorimeters in addition to the Endcap Calorimeters. This analysis therefore prioritizes geometrical acceptance.

The primary vertex of the event is defined as the one with the largest $\sum p_T^2$ where the sum is over all tracks in the ID with $p_T > 0.4$ GeV, and must have at least 3 such tracks. Combined muons, reconstructed in both the ID and the MS, are selected where possible, with segment-tagged muons and MS standalone muons accepted where necessary\footnote{The different muon reconstruction types are explained in Section 5.1} to extend the acceptance to the full range $|\eta| < 2.7$. Electron candidates are identified as clustered energy depositions in the Electromagnetic Calorimeter satisfying a set of requirements \cite{24} on the longitudinal and transverse shower shapes, associated with a well-reconstructed track in the ID. Both charged leptons are required to be isolated, and their tracks are required to be close to the primary vertex. The leading charged lepton candidate in the event is required to have $p_T > 22$ GeV, and the sub-leading one is required to have $p_T > 10$ GeV.

Jets are measured as topological clusters \cite{122} in the Hadronic Calorimeters, reconstructed using the anti-$k_t$ algorithm \cite{60} with a distance parameter $R = 0.4$. The measured transverse energy of jets, corrected to remove contributions from pileup events, is required to satisfy $E_T > 25$ GeV for $|\eta| < 2.4$, and $E_T > 30$ GeV for $2.4 < |\eta| < 4.5$. To further reduce contributions from pileup, if a jet has $E_T < 50$ GeV, at least 50% of the summed scalar $p_T$ of associated tracks must come from tracks associated with the primary vertex. Light-quark jets are identified from $b$-quark jets using a $b$-tagging algorithm \cite{20} to define a veto. The algorithm is based on the presence of displaced
vertices, the impact parameter of tracks in the jets and the reconstruction of heavy-flavour hadronic
decays when possible. The light-quark jet rejection factor, defined as the inverse of the rate for
mistaking a light-quark jet for a $b$-quark jet, is shown as a function of the $b$-tagging efficiency for
different $b$-tagging algorithms in Figure A.6. The chosen working point is at a $b$-tagging efficiency
of 85%, corresponding to a light-quark jet rejection factor of 10.

Missing transverse momentum \([21, 22, 28]\) measured using calorimeter energy deposits was
discussed in Section 4.6. In addition, a track-based $p^{\text{miss}}_T,\text{track}$ is defined using only the momentum of
tracks as measured from the ID and the MS.

![Figure A.6: Light-quark jet rejection factor as a function of the $b$-tagging efficiency for differ-
ent $b$-tagging algorithms [20]. The algorithm MV1, used in this analysis, is based on
a neural network taking the output weights from the IP3D, SV1 and JetFitterCombNN
algorithms as inputs; these algorithms are described in Ref. [17]. SV0, a predecessor
to SV1, is described in Ref. [16]. The JetFitterCombNNc algorithm makes use of the
JetFitterCombNN neural network trained to reject $c$-jets instead of light-quark jets [20].](image)

A.2.2 Backgrounds

Many physical processes can produce the same final state as the signal. The main background to this
search comes from top-quark backgrounds, namely $t\bar{t}$ and single-top production, where one or more
$b$-quarks from the decay is mis-identified as light-quark jet. Diboson backgrounds from $WW$, $WZ$, $W\gamma^*$(*) and $ZZ$ production in association with at least two jets in the initial state are also important,
as well as gluon-fusion Higgs and $Z/\gamma^*$ production in association with at least two jets. Finally,
backgrounds where at least one jet is mis-identified as a charged lepton, namely $W +\text{jets}$ and QCD multi-jets, also contribute to the event expectation in the Signal Region (SR). They are evaluated from data.

The main challenge of this analysis is therefore to effectively separate the small signal from background processes with relatively large cross sections. This is most easily seen when comparing the cross section times branching fraction of the VBF Higgs signal in the $\ell\nu\ell\nu$ channel (with $M_H = 125$ GeV):

$$\sigma(\text{VBF} H) \cdot B(H \rightarrow WW^* \rightarrow \ell\nu\ell\nu) \approx (1.58 \text{ pb})(1.0\%) = 15.8 \text{ fb} \quad (A.3)$$

which for an integrated luminosity of 20.7 $\text{fb}^{-1}$ corresponds to a total expected number of only 327 events, to the significantly larger cross sections for background processes illustrated in Figure 2.8.

### A.2.3 Topology of VBF $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ Events

Fortunately, the VBF $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ signal has a unique topology which makes the search possible. First, as illustrated in the Feynman diagram of Figure A.4, two back-to-back, forward, light-quark jets are expected. As mentioned in Section A.2.1, a $b$-tag veto is imposed with an efficiency of 85% to reject $b$-quark jets measured within the acceptance of the Inner Detector, in order to reject top-quark backgrounds.

The jets in signal are further expected to have a large opening angle in rapidity $\Delta y_{jj}$ and a high dijet invariant mass $m_{jj}$. Vetoing events with additional jets between the two leading jets also helps to reject background: this requirement is called the “central-jet veto”. Additionally, both jets are expected to be more forward than the leptons from the $W$ decays. Therefore, it is possible to define a variable called “lepton $\eta$-centrality”:

$$C_\eta = \sum_\ell \left| \frac{\sum \Delta \eta_{\ell j}}{\Delta \eta_{jj}} \right|$$

(A.4)

where the sums are over the two leading charged lepton candidates and the two leading jets. The quantity inside the absolute value is equal to unity when the lepton is collinear with one of the jets, smaller when it is between the two jets, and larger otherwise. In top-quark decays, since one lepton and one (mis-tagged) jet come from each top quark, the separation between each lepton and the closest jet is expected to be smaller than in signal; this is also captured by lepton $\eta$-centrality.

Another powerful variable against the $t\bar{t}$ background is $\sum m_{ij}$, where the sum is over all four pairs made from the two leading charged lepton candidates and the two leading jets. In $t\bar{t}$ events, $m_{ij} < M_t$ for two of these pairs, so $\sum m_{ij}$ is in general smaller for this background with respect to signal. The variable $m_T$ introduced in Equation A.2 also discriminates against $t\bar{t}$ and other backgrounds. In what follows, $p_{T,\text{track}}^{\text{miss}}$ is used in place of $E_T^{\text{miss}}$ in the definition of $m_T$ to improve the resolution.
of this observable. Background events with top quarks also tend to have a larger total transverse momentum, defined as the following vectorial sum:

$$ p_T^{\text{tot}} = \sum_\ell p_T^{\ell} + \sum_j p_T^{j} + p_T^{\text{miss,track}} $$

(A.5)

This variable is essentially a measure of the soft activity from initial-state radiation.

Backgrounds from $Z/\gamma^*$ production are also important: decays to $\tau^+\tau^-$ where both $\tau$ leptons decay leptonically potentially affect both search channels, while $Z/\gamma^* \rightarrow e^+e^-$ and $Z/\gamma^* \rightarrow \mu^+\mu^-$ only affect the SF channel. The photon and $Z$ peak regions are vetoed using window cuts on the invariant masses $m_{\ell\ell}$ and $m_{\tau\tau}$, where $m_{\tau\tau}$ is computed under the assumption that neutrinos from $\tau$ decays are collinear with their associated charged lepton. As well, since there is no real missing transverse momentum coming from $Z/\gamma^*$ production in the SF channel, requiring large values of $E_T^{\text{miss}}$ and $p_T^{\text{miss,track}}$ helps to reduce this background.

Finally, signal events have two leptons of opposite electric charge, going preferentially in the same direction due to spin correlations in the $H \rightarrow WW^*$ decay. Indeed, since the Higgs boson has spin zero, the two $W$ bosons must have opposite spin. Then, the $W^-$ decays to $\ell^-\bar{\nu}$, and the $W^+$ to $\ell^+\nu$. But antineutrinos always have right-handed helicity, i.e. their spin along the direction of motion is oriented in the same direction as their velocity, while neutrinos always have left-handed helicity. It follows, as illustrated in Figure A.7, that following a $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ decay, if the $W$ bosons have non-zero spin in the direction of motion, then the charged leptons travel in the same direction, while the neutrinos travel in the other. As a result, the charged lepton candidates are expected to have a small azimuthal separation $\Delta\phi_{\ell\ell}$ and a low dilepton invariant mass $m_{\ell\ell}$.

Figure A.8 shows a candidate signal event found in the dataset collected in 2012, which displays all the characteristics described above.

---

**Figure A.7**: Possible spin configurations following a $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ decay where the $W$ bosons have non-zero spin in the direction of motion. The thin arrows represent velocity, and the thick arrows represent spin. The charged leptons preferentially travel in the same direction. Figure credit: K. van Nieuwkoop.
Figure A.8: VBF $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ candidate event. The angular separation between the muon candidate (orange) and the electron candidate (green) is small, with $E_T^{\text{miss}}$ (red) opposite them. The two leading jets (blue) are very forward [33].
A.2.4 Boosted Decision Trees

The first multivariate analysis technique used in this analysis is based on BDTs, which constitute a machine-learning algorithm aiming to classify events as signal or background based on selected discriminating variables. The implementation is based on the Toolkit for Multivariate Data Analysis (TMVA) environment [118].

As a first step in any machine-learning algorithm, the sample of simulated background and signal events is separated into two halves. The first half, called the “training sample”, is used to train the algorithm to discriminate signal from background, while the other is an independent “test sample”, used to quantify the performance of the algorithm and provide background and signal predictions in the SR. In practice, two different BDTs are trained: one on the first half of the sample, to apply on the second, and vice versa for the other. This procedure, called cross-evaluation, allows to keep using the entire simulated sample to estimate the expected signal and background in the SR.

In order to understand how a BDT algorithm differs from a traditional cut-based approach to data analysis, it is helpful to first consider a single decision tree. Decision trees consist of a succession of cuts over the discriminating variables: an example is shown in Figure A.9. First, starting from a “root node” containing all the events used for training, a scan is performed to find the cut that best separates signal from background. Each side of this cut then defines a sub-node, which is further split. The process continues from each sub-node as long as the number of events in a node is above a minimum value; otherwise, the node ceases to be split and is called a “leaf node”. Each leaf node is identified as a “signal node” or a “background node” based on its contents.

Figure A.9: Schematic view of a simple decision tree, used to separate signal (S, blue) from background (B, red) [118].
Figure A.10 illustrates how using a single decision tree already constitutes an improvement over a simple cut-based analysis. For an example distribution of training background and signal events, the decision tree does find the region that would be selected using successive cuts, but also finds other signal-rich regions that would otherwise have not been identified. The importance of setting a minimal number of events per leaf node is also illustrated: not having this requirement can result in overtraining, whereby unphysical signal-rich regions are found that are specific to the training sample. This would result in an unreliable decision tree that performs very well, by construction, on the sample used to train it, but not as well on statistically independent samples.

To further improve the discriminating power and stability of this method, the initial decision tree is then “boosted”. First, the events from the training sample that were mis-classified in the decision tree are identified, and these events are given a higher weight. A second decision tree is then trained from this re-weighted sample: by construction it will improve the classification of the events with larger weights. This process is repeated until hundreds of decision trees are trained. The collection of all the decision trees thus trained is the Boosted Decision Tree.
Finally, for each decision tree in the BDT, each event is given a score of $-1$ if it ended in a background node, and $+1$ if it ended in a signal node. The final discriminant, called the “BDT score”, is defined for each event as the weighted average of the scores from all trees. The information from the discriminating variables and correlations between them allowing to separate signal from background is thus summarized by a single value assigned to each event.

### A.2.5 Matrix Element Method

A complementary approach to machine-learning algorithms such as BDTs is the Matrix Element (ME) method, an elegant approach from first principles which aims to assign to each event a probability to have been produced by a given physical process. Following a given physical process $i$, final-state particles with four-momenta $y$ are produced with probability $P_i(y)$. The probability to observe a set of physical observables $x$ in the detector is then given by

$$P_i(x) = \int P_i(y) T(x,y) dy$$

where the transfer functions $T(x,y)$ relate the particle-level quantities to the physical observables. The transfer functions used are often Dirac $\delta$-functions: this is justified by the good resolution of the detector to measure the angular coordinates and the momentum of charged lepton candidates, as well as the angular coordinates of jets. On the other hand, it is important to use Gaussian or even more involved distributions for the energy of jets, to account for possible differences between measured and particle-level values due to the detector resolution and thereby improve the sensitivity of the method.

Given the energy of collisions at the LHC, the mass of all final-state particles can safely be neglected. In the search for VBF $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$, the physical observables $x$ are therefore the three-momenta of the two charged lepton candidates and of the two leading jets, and the two components of $E_T^{\text{miss}}$. In contrast, $y$ represents the corresponding observable particle-level quantities: the three-momenta of the four leptons and of the two leading quarks or gluons in the final state. Four out of the six components of the neutrino three-momenta are unobservable, and are assigned $T(x,y) = 1$.

In the context of a proton-proton collision with colliding partons $q_1$ and $q_2$, Equation (A.6) becomes

$$P_i(x) = \frac{1}{\sigma_i} \int f(q_1) f(q_2) |\mathcal{M}_i(y)|^2 T(x,y) d\Phi(y) dq_1 dq_2$$

In this equation, the cross section $\sigma_i$ of the process of interest is used to normalize the result, ensuring that $P_i(x)$ is a probability density. The integrand includes the Parton Distribution Functions $f(q)$, which depend on the four-momenta $q$ of initial-state partons, the square of the matrix element $\mathcal{M}_i(y)$, from which the method gets its name, and $d\Phi(y)$ represents the phase space over which the integration is taken.
From the event-by-event probabilities calculated for signal and background processes, it is possible to define the Event Probability Discriminant (EPD):

\[
\text{EPD} = \frac{P_{\text{sig}}}{P_{\text{sig}} + \sum_{i \in \text{bkg}} \alpha_i P_i}
\]

(A.8)

where the coefficients \( \alpha_i \) are optimized to find the EPD definition with the best discriminating power. In theory, these coefficients should be close to unity; in practice they can vary due to the fact that the acceptance and efficiency for distinct background sources differ.

### A.3 Contributions to the Analysis

This section documents the main areas where contributions were made to this analysis in parallel with the work presented in the main body of this dissertation.

#### A.3.1 Statistical Framework

The implementation of the statistical analysis was conducted in parallel with the optimization and background estimation procedures described below, within an existing framework created for earlier versions of the analysis. Figure A.11 shows a schematic diagram of the analysis fit model. As indicated above, the search is divided into two channels, depending on the flavour of the leading charged leptons: in the Different-Flavour (DF) channel, these leptons are one muon and one electron, while in the Same-Flavour (SF) channel, they are either two electrons or two muons.

All parameters are fitted simultaneously. Three Control Regions (CRs) are defined in order to constrain the corresponding backgrounds using data:

- Top CRs, defined by requiring exactly one \( b \)-tagged jet in the event, are used in both channels to correct the normalization and shape of the predictions from MC simulations of background processes involving top quarks. The CRs are distinct for the DF and SF channels in order to benefit from the higher purity in the DF sample.

- The \( Z/\gamma^* \rightarrow \tau^+ \tau^- \) CR, discussed in Section A.3.6, is used to normalize the yields predicted by simulation of this background to those observed in data.

- The \( Z/\gamma^* \rightarrow \ell^+ \ell^- \) CR, discussed in Section A.3.7, is used in the SF channel to obtain a fully data-driven estimate of this background.

The parameters used to correct the simulations are called Normalization Factors (NFs), defined as the ratio \( N_{\text{data,CR}}/N_{\text{MC,CR}} \). The expected number of events in the SR is then given by

\[
N_{\text{exp,SR}} = \frac{N_{\text{data,CR}}}{N_{\text{MC,CR}}} \cdot N_{\text{MC,SR}} = \frac{N_{\text{MC,SR}}}{N_{\text{MC,CR}}} \cdot N_{\text{data,CR}}
\]

(A.9)
where the ratio in the second expression is commonly called the “extrapolation factor”.

The systematic uncertainties are implemented as nuisance parameters which are then marginalized. The value for the VBF signal strength is extracted from the fit, along with the signal significance. All optimizations below seek to maximize the expected significance, calculated from the estimated signal and background yields in the fit regions.

**Figure A.11:** Schematic diagram of the analysis fit model.

### A.3.2 Event Pre-Selection Optimization

In addition to the selection of physical objects described in Section A.2.1 and the $b$-tag veto, a few more criteria are applied to the events before applying the multivariate analysis techniques:

- The requirement $m_{\tau\tau} < 66$ GeV, called the $Z/\gamma^* \rightarrow \tau^+\tau^-$ veto, is applied in both channels to allow the definition of the $Z/\gamma^* \rightarrow \tau^+\tau^-$ CR.

- The photon peak is absent from the $Z/\gamma^*$ background simulation; this is replicated in the analysis by requiring $m_{\ell\ell} > 10$ GeV in the DF channel and $m_{\ell\ell} > 12$ GeV in the SF channel. The cut in the SF channel also removes potential background contributions from $J/\psi$ and $\Upsilon$ meson decays.

- The region of the $Z$ peak and above is set aside in the SF channel for use in the data-driven estimate of the $Z/\gamma^* \rightarrow \ell^+\ell^-$ background; the same is true of the low-$E_T^{\text{miss}}$ region. The requirements $m_{\ell\ell} < 75$ GeV and $E_T^{\text{miss}} > 45$ GeV are therefore applied in this channel.
Further, it was found during the BDT optimization described in the next section that the same BDT performs well in both the DF and SF channels, as long as additional cuts are applied to reduce the background from $Z/\gamma^* \rightarrow \ell^+ \ell^-$. This greatly simplifies the analysis, as the systematic uncertainties and NFs for most processes are then the same in both channels.

The original cuts designed for this purpose were on the transverse momentum $p_{T,\ell\ell\text{jets}}$ of the system consisting of all the lepton candidates and jets in the event, and the soft hadronic recoil fraction $f_{\text{recoil}}$, defined as follows:

\[ p_{T,\ell\ell\text{jets}} = \left| \sum_{\ell} \vec{p}_{T,\ell} + \sum_{j} \vec{p}_{T,j} \right| \]  

\[ f_{\text{recoil}} = \frac{\left| \sum_{\text{opp. jets}} \vec{p}_{T,j} \right|}{p_{T,\ell\ell\text{jets}}} \]  

where in the definition of $f_{\text{recoil}}$, the sum is over the jets opposite to the system of leptons and jets, i.e. over the jets with azimuthal separation from $\vec{p}_{T,\ell\ell\text{jets}}$ larger than $3\pi/4$. These variables discriminate against the $Z/\gamma^*$ background because of the absence of neutrinos in this background’s final state, resulting in smaller values of $p_{T,\ell\ell\text{jets}}$ and larger values of $f_{\text{recoil}}$. A good operating point consists of the requirements $p_{T,\ell\ell\text{jets}} > 25$ GeV and $f_{\text{recoil}} < 0.2$.

On the other hand, it was later found that replacing these cuts by a more straightforward requirement on $p_{T,\text{track}}^{\text{miss}} > 40$ GeV gives the same performance, with less reliance on the modelling of soft jets. This later requirement was therefore adopted for the analysis. While any of these variables could have been included in a dedicated BDT for the SF channel, the advantages of using the same BDT in both channels are more important than the marginal increase in sensitivity that such an approach would bring.

### A.3.3 Boosted Decision Tree Optimization

Following the event pre-selection, BDTs are trained and applied to separate background from signal. The identity of the input variables used in the BDTs is determined using the algorithm illustrated in Figure A.12, aiming to find the simplest BDT configuration that still maximizes performance.

First, an initial BDT is trained with a large number $N$ of input variables, and the performance of this BDT is evaluated on the test samples using the expected significance as figure of merit. This value is considered the performance ceiling. The initial BDT is then simplified by considering $N$ different BDT, each with one less input variable than the initial BDT. The input variable missing from the most performant of these simplified BDTs is the one that provides the least marginal improvement. If the performance of this best simplified BDT is similar to that of the initial BDT, the algorithm is repeated using the best simplified BDT as the new initial one.

It is following this procedure that eight variables described in Section A.2.3 were identified as the best input variables for the BDT: $m_T$, $m_{\ell\ell}$, $\Delta \phi_{\ell\ell}$, $m_{jj}$, $\Delta y_{jj}$, $p_{T,\ell\ell\text{jets}}^{\text{tot}}$, $\sum m_{\ell j}$ and $C_\eta$.  

187
The BDT configuration settings are also optimized, via a grid scan of the different possibilities. The optimal and most stable configuration is found to be a BDT made of 1000 trees, each with a maximum depth of 5 levels and a minimum of 1000 events in each leaf node, with additional settings related to the boosting algorithm also optimized.

Finally, the binning of the BDT score used as the final discriminant is also optimized. The binning should be as fine as possible, to limit the loss of information, while being coarse enough to make possible a reliable estimate of expected signal and background yields as a function of the BDT score. The problem being that the importance of statistical errors grows when finer bin sizes are used.

A grid scan was originally used to find the optimal bin boundaries with the expected significance as the figure of merit, however because of the need to repeat this step for each BDT tried (as opposed to the configuration settings which remain the same) the following, faster method is used instead. The bin boundaries are set by integrating the signal and background distributions from the high end of the BDT score distribution, to find the point maximizing the Poisson significance

\[
Z_{\text{Poisson}} = \sqrt{2(N_{\text{sig}} + N_{\text{bkg}}) \cdot \ln(1 + N_{\text{sig}}/N_{\text{bkg}})} - 2N_{\text{sig}}
\]

(A.12)

This figure of merit is a closer approximation to the actual significance than the Gaussian approximation \(N_{\text{sig}}/\sqrt{N_{\text{bkg}}},\) which is only valid for \(N_{\text{sig}} \ll N_{\text{bkg}}\) [84]. When a boundary is set, the integration restarts from this point to find the next bin boundary. Three bin boundaries are typically set by the algorithm, defining four bins. The three most signal-rich bins form the SR, and the background-rich bin with events having the lowest BDT scores is used as a Validation Region (VR), to verify the modelling of input variables and correlations.

A.3.4 Modelling Studies

Since the BDT is trained on MC simulated events, it is essential for the method to be valid that the modelling of data by the simulation is correct. The BDT score itself is verified in both search...
channels in the Top CR, and these histograms are shown in Figure A.13. The BDT training variables are also verified in both search channels, in the Top CR as well as in the low-BDT VR. These histograms are shown in Figures A.14, A.15, A.16 and A.17. The agreement between data and the sum of expected backgrounds is satisfactory in all the distributions.

Figure A.13: Histograms of the BDT score in the Top CR for the DF channel (top) and the SF channel (bottom). The uncertainties shown on the background estimate are statistical only.
Figure A.14: Histograms of the BDT training variables in the Top CR for the DF channel. The uncertainties shown on the background estimate are statistical only.
Figure A.15: Histograms of the BDT training variables in the Top CR for the SF channel. The uncertainties shown on the background estimate are statistical only.
Figure A.16: Histograms of the BDT training variables in the low-BDT VR for the DF channel. The uncertainties shown on the background estimate are statistical only.
Figure A.17: Histograms of the BDT training variables in the low-BDT VR for the SF channel. The uncertainties shown on the background estimate are statistical only.
A.3.5 Correlation Studies

In addition to the BDT discriminating variables themselves, it is very important to verify that the correlations between these variables are well-modelled in the simulation, in order to guarantee the validity of the conclusions to be drawn from the BDT score distribution. Two ways to make this verification are employed.

The first way is based on the sample approximation to the Pearson correlation coefficient $\rho$:

\[
\rho(x,y) = \frac{(x - \bar{x})(y - \bar{y})}{\sqrt{\sum (x - \bar{x})^2 \sum (y - \bar{y})^2}}
\]

where $\bar{x}$ and $\bar{y}$ are the sample means of the distributions of the variables $x$ and $y$ respectively, and the sums in the denominator are taken over either data or the expected backgrounds. For each pair of variables used in the BDT, the distribution of $\rho(x,y)$ is compared between data and the sum of expected backgrounds. An example of such a comparison is shown in Figure A.18. The agreement is reasonable for all pairs of variables.

![Figure A.18: Distribution of $\rho(m_{\ell\ell}, m_T)$ in the low-BDT VR. The mean and its statistical uncertainty is indicated for background and data.](image)

As a more quantitative improvement on this method, it is also possible to monitor projections of the N-dimensional phase space onto 2D profile histograms. Figure A.19 shows an example of this test: the mean value of one variable is shown in bins of the other, and vice versa. This test is repeated for all pairs of variables, as well as between all variables and the BDT score. $\chi^2$-probabilities are
calculated to quantify the agreement between data and the sum of expected backgrounds. In the low-BDT VR, the $\chi^2$-probability is larger than 0.3% for all histograms but one ($C_\eta$ vs. $m_{\ell\ell}$), and larger than 5% for two others out of 72. In these three cases the deviations are evenly distributed on each side of a unit ratio, and are therefore ascribed to statistical fluctuations.

Figure A.19: Example pair of 2D correlation histograms, for data (black) compared with the sum of expected backgrounds (red). The green colour indicates that the $\chi^2$-probability is larger than 5%.

A.3.6 Background Estimation and Systematic Uncertainties for $Z/\gamma^* \rightarrow \tau^+ \tau^-$

As mentioned in Section A.3.1, a $Z/\gamma^* \rightarrow \tau^+ \tau^-$ CR is used to estimate a Normalization Factor (NF) for this background in the DF channel. The $Z/\gamma^* \rightarrow \tau^+ \tau^-$ CR is defined in both the DF and SF channels as follows:

- All pre-selection cuts except the $Z/\gamma^* \rightarrow \tau^+ \tau^-$ veto;
- $m_{\ell\ell} < 80$ GeV in the DF channel (the SF channel pre-selection includes $m_{\ell\ell} < 75$ GeV);
- $|m_{\tau\tau} - M_Z| < 25$ GeV;
- BDT score requirement corresponding to the SR definition.

The cuts on $m_{\ell\ell}$ and $m_{\tau\tau}$ increase the $Z/\gamma^* \rightarrow \tau^+ \tau^-$ purity of the CR. Figure A.20 shows the $m_{\ell\ell}$ and $m_{\tau\tau}$ distributions in the $Z/\gamma^* \rightarrow \tau^+ \tau^-$ CR for the combination of the DF and SF channels.

The resulting NF, derived from this combined CR, is $0.9 \pm 0.3$ (statistical uncertainty only), which is compatible with unity. It is applied in the fit, and its uncertainty is considered as uncorrelated between the different bins of the SR. Despite the fact that this statistical uncertainty is relatively large due to the small event yield in the CR, its impact on the result is negligible because the contribution of the $Z/\gamma^* \rightarrow \tau^+ \tau^-$ process to the total background in the DF SR is small in comparison to the ones from dominant sources such as $t\bar{t}$ production.
Figure A.20: Distributions of $m_{\ell\ell}$ (left) and $m_{\tau\tau}$ (right) in the $Z/\gamma^* \rightarrow \tau^+\tau^-$ CR for the combination of the DF and SF channels.

A.3.7 Background Estimation and Systematic Uncertainties for $Z/\gamma^* \rightarrow e^+e^-$, $\mu^+\mu^-$ in the Same-Flavour Channel

To estimate the $Z/\gamma^*$ background in the SR of the SF channel, the BDT analysis makes use of a fully data-driven method inspired from ABCD techniques. In these techniques, the background estimate in one region of phase space, called region A (typically the SR), comes from three neighbouring regions. The shape of the background in terms of the variable of interest is taken from one of these neighbouring regions, called region B, and the extrapolation factor going to region A from region B is taken from the other two regions, called C and D.

Here, the $Z/\gamma^*$ shape as a function of BDT score is taken from the data in a low-$E_T^{miss}$ $Z$ CR ($25 \text{ GeV} < E_T^{miss} < 45 \text{ GeV}$, using calorimeter-based $E_T^{miss}$), corrected by subtracting the other background contributions in this region. The $Z/\gamma^*$ estimate is extrapolated to the SR using the $E_T^{miss}$ cut efficiency calculated from data on the $Z$ peak. Table A.1 illustrates the regions used in this method: the indicated cuts are used in addition to the rest of the selection in the SF channel.

The $Z/\gamma^*$ estimate in each BDT score bin $i$ is then determined by:

$$N_{Z/\gamma}^{SR,i} = N_{Z/\gamma}^B \cdot \frac{N_{Z/\gamma}^C}{N_{Z/\gamma}^D}$$

(A.14)

where on the right-hand side $N_{Z/\gamma} = (N_{data} - N_{non-Z/\gamma MC})$. The resulting NFs are compatible with unity in each bin of BDT score. Figure A.21 shows the BDT score distributions from each region.

The $E_T^{miss}$ cut efficiency measured in data from regions C and D is $0.43 \pm 0.03$, consistent with the value of $0.47 \pm 0.04$ measured from $Z/\gamma^*$ MC simulation. The statistical uncertainty on the efficiency measured from data is propagated to the fit, fully correlated across bins of BDT score.
Table A.1: Summary of the region definitions for the $Z/\gamma^{*}$ estimation technique used in the SF channel of the analysis.

<table>
<thead>
<tr>
<th>Region A (SR)</th>
<th>Region C</th>
</tr>
</thead>
<tbody>
<tr>
<td>$E_{T}^{\text{miss}} &gt; 45$ GeV</td>
<td>$E_{T}^{\text{miss}} &gt; 45$ GeV</td>
</tr>
<tr>
<td>$m_{\ell\ell} &lt; 75$ GeV</td>
<td>$</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>Region B (Z CR)</th>
<th>Region D</th>
</tr>
</thead>
<tbody>
<tr>
<td>$25$ GeV &lt; $E_{T}^{\text{miss}}$ &lt; $45$ GeV</td>
<td>$25$ GeV &lt; $E_{T}^{\text{miss}}$ &lt; $45$ GeV</td>
</tr>
<tr>
<td>$m_{\ell\ell} &lt; 75$ GeV</td>
<td>$</td>
</tr>
</tbody>
</table>

It is important to note from Equation A.14 that the $E_{T}^{\text{miss}}$ cut efficiency is not calculated in bins of BDT score, but rather obtained from the full sample after the selection: the $E_{T}^{\text{miss}}$ cut efficiency is assumed to be the same across the BDT score spectrum. This is necessary because little to no events from the Z peak populate the bins at high BDT score. This method therefore relies on the absence of correlation between the BDT score and $E_{T}^{\text{miss}}$. This fact is demonstrated by the similarity of the BDT score shapes from $Z/\gamma^{*}$ MC simulation in the SR and the Z CR, shown in Figure A.22. Even though the two shapes are compatible with each other within statistical uncertainties, the deviation from unity of the ratio between the two shapes in each bin is taken as an uncorrelated systematic uncertainty due to the non-correlation assumption.

A second assumption that has to be satisfied for the method to be valid is the absence of correlation between $E_{T}^{\text{miss}}$ and $m_{\ell\ell}$, which is necessary to extrapolate the $E_{T}^{\text{miss}}$ cut efficiency value from regions C and D to regions A and B. A closure test in $Z/\gamma^{*}$ simulation is performed to test this assumption, with the non-closure between the four regions quantified using a double ratio:

$$\text{non-closure} = \frac{N_{A,\text{high-BDT}}^{Z/\gamma^{*}}}{N_{B,\text{high-BDT}}^{Z/\gamma^{*}}} \times \frac{N_{C}^{Z/\gamma^{*}}}{N_{D}^{Z/\gamma^{*}}}$$  \hspace{1cm} (A.15)

where “high-BDT” corresponds to the BDT score bins used in the fit. Closure is valid within statistical uncertainty, with a non-closure value of $0.87 \pm 0.22$. The deviation of this value from unity is taken as a flat systematic uncertainty on the $Z/\gamma^{*}$ estimate in the SF channel.
Figure A.21: Distributions of BDT score in the SF channel, for the low-\(E_T^{\text{miss}}\) Z CR (top), region C (bottom left) and region D (bottom right).
Figure A.22: Illustration of the absence of correlation between the BDT score and $E^\text{miss}_T$ in simulated events. Specifically, the BDT score distribution shape, normalized to unity, is shown for $Z/\gamma^*$ simulation for the SR selection (red) and in the $Z$ CR selection (green). The rightmost three bins correspond to the fit SR, while the other two correspond to the low-BDT VR. The shapes are compatible within the statistical uncertainties.

A.3.8 Effect of Object Systematic Uncertainties on the BDT Score Shape

The systematic uncertainties affecting this search form two main categories: theoretical systematic uncertainties originating from the background and signal estimation techniques, and experimental systematic uncertainties coming from the measurement and identification of the physical final state objects. The latter category includes, in decreasing order of importance, the uncertainties on the energy resolution and scale of jets and electron candidates, the scale and resolution of missing transverse momentum, the muon momentum scale and resolution, the jet $b$-tagging efficiencies, the lepton reconstruction and isolation efficiencies, and others.

For each object systematic uncertainty source, dedicated signal and background samples are produced with the same simulated events as the nominal samples, but with the object property affected by the systematic uncertainty in question varied by the uncertainty range at 68% CL. The event selection is then applied on these samples, and a nuisance parameter is assigned to the shape differences in the discriminating variable between each varied sample and the nominal sample. An example object systematic uncertainty variation is shown in Figure A.23.
Figure A.23: Example object systematic uncertainty, due to the $p_T$ track resolution, on the BDT shape in the SR of the DF channel, for signal (left) and the sum of all backgrounds (right). The bottom inset shows the ratio between the varied and nominal samples. The yellow band represents the statistical uncertainty.

A.3.9 Comparison of Analysis Techniques

The new expected result based on BDT can be compared with the cut-based result to quantify the compatibility of the two approaches. As a first step, the overlap between the events in the SRs of the two analyses is quantified in MC simulation: the result is that 99% of the events in the SR of the cut-based approach are selected in the BDT SR. On the other hand, the signal acceptance of the BDT SR is more than double that of the cut-based SR.

Even then, the question of the compatibility of observed results remains: as it may turn out that the measured signal strength and significance will be very different between the two approaches, it is useful to quantify the expected range of differences between the two results. This question is addressed using pseudo-experiments drawn from simulated events.

In order to meaningfully compare the outcomes from the two analysis approaches, it is necessary to use the same pseudo-datasets in both cases. This forbids the method of drawing pseudo-experiments directly from the distributions of the discriminating variables, as was done for example is Chapter 8 to estimate expected limits and calculate $p$-values.

Instead, actual pseudo-datasets consisting of simulated events are drawn randomly from the MC samples used to evaluate the background and signal yields. The number of events drawn from each sample is itself drawn from a Poisson distribution with a mean corresponding to the expected yield from the given process following the event selection. Since the MC simulated events have different weights, the sampling method must take this into account, in that the probability to draw a given
event must be proportional to its weight. The following algorithm [102], illustrated in Figure A.24, is used for this purpose:

- For each sample, make a distribution of the cumulative event weight sum.
- Draw a random number between zero and the total weight sum: this determines which particular simulated event is picked-up.
- Repeat until the desired number of events is drawn.

Figure A.24: Illustration of the sampling algorithm to generate pseudo-experiments [102].

The method is applied here in the DF channel of the search, with a nominal signal strength corresponding to Standard Model predictions (\( \mu \equiv \sigma / \sigma_{SM} = 1 \)) injected in the pseudo-datasets. Two hundred pseudo-experiments are generated. The pseudo-datasets are treated like observed data samples in the statistical framework: for each of them in turn, the event selection is applied and the fits are run to obtain “observed” signal strength and significance values in each analysis method.

First, as a validation step, histograms of the number of events passing the selection are produced: the results are shown in Figure A.25 for the selection taking place in each analysis. As expected, the distributions obtained from pseudo-datasets follow the Poisson distributions corresponding in each case to the injected signal strength. This constitutes a non-trivial test that the pseudo-data are representative: while the distributions of the number of events in each pseudo-dataset match the Poisson distribution by construction before the selection, this is not necessarily the case after the event selection is applied.

The observed signal strength \( \hat{\mu} \) and observed significance from pseudo-experiments are shown
Figure A.25: Number of events in pseudo-datasets after the pre-selection of the BDT analysis (left) and the selection of the cut-based analysis (right).

in Figure A.26 for the BDT analysis, and in Figure A.27 for the cut-based analysis. On these histograms, the median expected significance from each analysis is also shown, and the median and average observed significances are shown to agree with these values. The agreement is also good between the median and average observed signal strength and the injected value \( \mu = 1 \).

One possibly surprising feature of these distributions is their width: for example, the BDT analysis has a percent-level chance to observe a significance as large as 4\( \sigma \) or, on the other side of the spectrum, not to observe any signal. For the cut-based analysis, this is explained by the width of the Poisson distributions shown in Figure A.25 (right): the low tail of the distribution with signal is around the median of the distribution without signal, while the high tail is significantly outside. This fact, in turn, is explained by the low expected signal yields in the SR. A similar argument would hold for the distribution in the BDT analysis after a cut on the BDT score.

For each pseudo-experiment, one can then directly compare the results obtained with the two analyses: histograms of the differences in signal significance and signal strength are shown in Figure A.28. The median and average differences are consistent with zero, indicating the absence of bias between the two analyses, however the width of this distribution indicates that potentially large differences are expected between the two observed results. This is explained like the variation in results within each analysis, by the fact that the expected signal yields are small.

Figure A.29 shows 2D histograms of the signal significance and strength observed in individual pseudo-experiments in the two analyses. The Pearson correlation coefficients are 0.63 for the ob-

---

3 The significance values shown here represent the status of the two analyses as of October 2013, and the performance in both of them has improved since.
Figure A.26: Signal significance (left) and strength (right) observed in 200 pseudo-experiments, with a BDT score as the discriminating variable.

Figure A.27: Signal significance (left) and strength (right) observed in 200 pseudo-experiments, with $m_T$ as the discriminating variable.

served significance values, and 0.65 for the signal strength. In addition, linearity tests are performed on these distributions to ensure that the methods are unbiased with respect to each other: these are shown in Figure A.30. The response of the methods is shown to be linear in both cases.

Notwithstanding these encouraging results, more verifications could be carried out in order to demonstrate that the results are representative of the expected range of possible outcomes:

- Performing the representativity test in bins of the SRs in addition to the inclusive tests shown...
Figure A.28: Differences in signal significance (left) and strength (right) observed in individual pseudo-experiments between the two analyses.

here would ensure that there are little shape distortions induced by the sampling algorithm.

- The signal strength distributions obtained with this event sampling algorithm could be compared, for individual analysis approaches, with the result of drawing pseudo-experiments directly from the respective distributions of the discriminating variables.

- Verifying the re-sampling distributions, i.e. histograms of the number of times each simulated event was drawn, would allow to quantify the maximum number of independent pseudo-experiments that can be drawn from the available simulated events.

This study will be repeated once the implementation of the re-analysis is complete, immediately before un-blinding the search.
Figure A.29: Two-dimensional histograms of the signal significance (left) and strength (right) observed in individual pseudo-experiments in the two analyses.

Figure A.30: Linearity test between the two analyses, for the signal significance (left) and signal strength (right).
A.3.10 Matrix Element Method Investigation

As introduced in Section A.2.5, at the beginning of this multivariate analysis effort two options were considered: the BDT method, and the ME method. For this analysis, the event-by-event probabilities from the ME method were calculated using two environments: MADWEIGHT [4, 10] (itself based on MADGRAPH [5]) for the $t\bar{t}$ and WW background hypotheses as well as for the VBF signal hypothesis, and independent software [153] for the $Z/\gamma^*$ and gluon-fusion Higgs background hypotheses. These probabilities were then combined into an EPD using Equation A.8, with the coefficients $\alpha_i$ optimized using a grid scan with the Poisson significance as the figure of merit.

Figure A.31 shows the modelling of the probability distributions in a previous definition of the Top CR in the DF channel. The agreement between data and the simulation is satisfactory. Figure A.32 shows the expected signal and background yields as a function of the event-by-event probabilities, illustrating the discriminating power of the method.

One of the main drawbacks of the ME method is the large computing times it requires. With integration times reaching up to a few minutes per event, the use of extensive computing resources was necessary, and the flexibility and usability of the method was therefore hampered.

Nevertheless, early results were very promising. The performance of both methods taken on their own were at first comparable, but it was observed that combining the EPD and the BDT score into a combined discriminant performed best. The best such combination found was the product EPD*BDT. The performance obtained with the combined discriminant was found to be up to 15% better than with either the EPD or the BDT score alone.

However, as the BDT optimization progressed, this additional gain started to fade, a plausible explanation being that the improvements to the BDTs were due to them making use of increasingly larger fractions of the available information relevant to discriminating signal from background. The disappearance of the gain from using a combined discriminant, along with the prohibitive computing times of the event-by-event probabilities, justified shelving this multivariate analysis method in favour of an analysis based solely on BDTs.
Figure A.31: Event-by-event probabilities calculated in a previous definition of the DF Top CR under the hypotheses of $t\bar{t}$ (top left), $WW$ (top right), gluon-fusion Higgs (centre left), $Z/\gamma^*$ (centre right) and VBF Higgs production (bottom). On the main histograms, the signal contribution is magnified by a factor 200 for visibility. The red line in the bottom inset indicates the bin-by-bin Poisson significance of the VBF signal.
Figure A.32: Event-by-event probabilities calculated in a previous definition of the blinded DF SR under the hypotheses of $t\bar{t}$ (top left), $WW$ (top right), gluon-fusion Higgs (centre left), $Z/\gamma^*$ (centre right) and VBF Higgs production (bottom). The red line in the bottom inset indicates the bin-by-bin Poisson significance of the VBF signal.
A.4 Conclusion and Outlook

In conclusion, the re-analysis of the Run-I dataset in the search for the process $H \rightarrow WW^* \rightarrow \ell\nu\ell\nu$ is well underway. Because the final state of interest is matched by a variety of diverse backgrounds, a multivariate analysis is found to be more powerful than a more traditional cut-based approach. Specifically, combined with refinements in the object selection, the new analysis strategy using BDTs is expected to bring an improvement of 40% in expected significance compared to the preliminary result.

Figure A.33 shows the preliminary ATLAS measurements from Run-I of the LHC of the signal strength ratios between the bosonic and fermionic Higgs production channels for the individual final states and their combination. All values are currently compatible with a unit ratio. Figure A.34 presents fit results for a parametrization of Higgs boson coupling strengths probing a universal scale factor for fermion couplings ($\kappa_F$) and another for boson couplings ($\kappa_V$). The measured couplings to fermions are compatible with predictions from the Standard Model, and although the couplings to bosons are more compatible with higher values, the measured values are consistent with the Standard Model prediction given the uncertainties.

Increasing the precision of Higgs boson coupling measurements might unveil a significant discrepancy, which could be indicative of physical processes beyond the Standard Model. Measuring the Higgs boson production and decay rates using LHC data from Run-II and beyond will complement direct searches for new particles, and is therefore an indispensable part of the high-energy physics programme for the next decades.
Figure A.33: Measurements of the signal strength ratios between the bosonic and fermionic Higgs production channels for the individual final states and their combination [34].

Figure A.34: Fit results for a parametrization of Higgs boson coupling strengths probing different scale factors for fermions and bosons [34].