UNVEILING NEW PHYSICS THROUGH ANGULAR DISTRIBUTIONS IN DILEPTON AND DIBOSON FINAL STATES AT THE LARGE HADRON COLLIDER

A Dissertation

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by

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Abstract

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This thesis focuses on searches for new physics effects using angular distributions in particle colliders. In particular, angular distributions in dilepton and diboson final states have been studied. Two type of new physics effects have been searched for: 1. Small deviations in dilepton distributions due to the presence of quantum interference effects involving dark matter particles, and 2. Signals from exotic heavy resonances in the rapidity distributions of $W\gamma$ production.

This document is structured as follows: Chapter 1 presents an introduction on the world of particle physics. This chapter is written in a language for the general public. Chapter 2 contains a small introduction to some aspects of particle physics phenomenology. This chapter is a little more technical than the previous one and it introduces some of the concepts needed for later chapters. Chapter 3 and 4 present the main results of this thesis using dilepton and diboson production at the LHC. The general conclusions are presented in Chapter 5. Finally, the two appendices present some of the publications that were made during the development of this thesis that do not fit into the general topic of the document.
DEDICATION

Just like my heart, this PhD is only half mine.
The other half belongs to my wife Mileidy Varela.
**CONTENTS**

Figures ................................................................. v

Tables ................................................................. ix

Acknowledgments ...................................................... x

Chapter 1: Introduction .............................................. 1
  1.1 A Brief Retrospect of Particle Physics ...................... 1
  1.2 Particle Physics Today ....................................... 6
  1.3 The Approach to Particle Physics in this Thesis ........... 10

Chapter 2: Elements of Particle Physics Phenomenology .......... 12
  2.1 Collider Physics .............................................. 12
  2.2 Dark Matter .................................................... 19

Chapter 3: Dark Matter Signatures in Dilepton Final States ...... 24
  3.1 Motivation ...................................................... 24
  3.2 Simplified Dark Matter Models .............................. 26
  3.3 Dark Matter Signatures at Colliders ....................... 29
  3.4 Constraints on the DM Parameter Space .................... 38

Chapter 4: New Physics Effects in Diboson Final States .......... 48
  4.1 Motivation ...................................................... 48
  4.2 Radiation Amplitude Zero .................................... 50
  4.3 New Physics Signatures in $W\gamma$ ........................ 56
  4.4 Results ......................................................... 58

Chapter 5: Summaries and Conclusions ............................. 63
  5.1 New Physics Signals in Dileptons ........................... 63
  5.2 New Physics Signals in Dibosons ............................ 64
  5.3 A General Conclusion ........................................ 65
<table>
<thead>
<tr>
<th>Appendix A: The Higgs mass in Supersymmetric Extensions of the SM</th>
<th>67</th>
</tr>
</thead>
<tbody>
<tr>
<td>A.1 Theoretical Set-up</td>
<td>67</td>
</tr>
<tr>
<td>A.2 Limits on the $Z'$ Mass</td>
<td>68</td>
</tr>
<tr>
<td>A.3 The Higgs Mass</td>
<td>69</td>
</tr>
<tr>
<td>A.4 Stops Masses</td>
<td>71</td>
</tr>
<tr>
<td>A.5 Conclusions</td>
<td>73</td>
</tr>
<tr>
<td>Appendix B: Leptoquarks and the R(D) Anomaly</td>
<td>74</td>
</tr>
<tr>
<td>B.1 Theoretical Framework</td>
<td>75</td>
</tr>
<tr>
<td>B.2 Operator Analysis</td>
<td>76</td>
</tr>
<tr>
<td>B.3 Leptoquark Interactions and Flavor Structure</td>
<td>78</td>
</tr>
<tr>
<td>B.4 Constraints and Results</td>
<td>81</td>
</tr>
<tr>
<td>B.5 Summary and Conclusions</td>
<td>86</td>
</tr>
<tr>
<td>Bibliography</td>
<td>88</td>
</tr>
</tbody>
</table>
FIGURES

1.1 The Standard Model of Particle Physics. Each box shows in the top-left corner the mass of each particle, its electric charge, and another property called spin. The mass is provided in units of energy according to Eq. 1.1. The quarks and gluons feel the strong force. Charged leptons, neutrinos, the Z and the W boson feel the weak force. All charged particles can feel the electromagnetic force. The Higgs particle interacts with all massive particles. (Figure from *Wikipedia Commons*.)

1.2 Illustration of the Higgs mechanism. The Higgs field (water) permeates all the universe. Each particle (swimmer) interacts differently with the Higgs field (number of floaters attached). The particles that interact the most with the Higgs field are more massive (the more floaters a swimmer has attached to himself, the slower he moves in the water).

2.1 Invariant mass of electron-positron pairs for 13 TeV and 35.9 fb⁻¹. (The label \(m(ee)\) does not refer to di-electron, but to electron-positron production.) Different colors represent the expected number of events from different processes as indicated by the colored boxes. The green line is a simulated signal from higher dimensional operators related to new physics (NP) effects. (Plot taken from [6].)

2.2 Collins-Soper distribution of \(e^+e^-\) events for 8 TeV and 20.3 fb⁻¹. Different colors represent the expected number of events from different processes as indicated by the colored boxes. The green line is a simulated signal from higher dimensional operators related to NP effects. (Plot taken from [8].)

2.3 Kinematic distributions for \(WZ\) production. The ATLAS collaboration shows the transverse momentum (left) and difference in rapidities of the lepton and the Z boson (right) for 13 TeV and 36.1 fb⁻¹. (Plot taken from [9].)

2.4 Dark Matter annihilation into neutrinos through a \(Z\) boson.

2.5 Dark Matter production at colliders. The actual signal at the LHC are *monophoton* events.

2.6 Searches of two jets + MET at the LHC. These can be interpreted as DM searches for models where DM (\(\tilde{\chi}^0_1\)) interacts with quarks through a hypothetical mediator particle (\(\tilde{q}\)). (Plot taken from [19].)
3.1 Feynman diagrams for dilepton production at the LHC. On top is
the Standard Model Drell-Yan process at tree level. The middle row
shows the box contributions from pseudo-Dirac DM with scalar medi-
ators. The bottom row shows the same from pseudo-complex DM
with fermion mediators. See the text and Table 3.1 for more details. 30

3.2 Dilepton invariant mass (left-top), CS angle (right-top), and angular
asymmetry distributions (bottom). The blue line represents the SM
background from Drell-Yan production. The orange, green, red, and
purple represent the pseudo-Dirac, pseudo-complex, Majorana, and
real scalar cases. Solid lines represent models with right-handed quarks
and right-handed leptons (RR) whereas dashed lines represent the RL
models. Signal lines are plotted at the benchmark point \( \lambda = 2.0, m_\chi = 500 \text{ GeV}, \) and \( m_\phi = 550 \text{ GeV}. \) 34

3.3 The bounds on the models at two different hierarchies between \( m_\phi \)
and \( m_\chi \). Bounds at 95\% CL from the LHC are obtained from mea-
urements at 8 TeV and 20 fb\(^{-1}\); the orange (blue) curves depict
dilepton bounds on the RR (RL) models, and are solid (dashed) for
\( m_{ee} (\cos \theta_{CS}) \) bounds; the purple regions are excluded by jets + MET
searches. The green curves are 90\% CL Xenon1T constraints on spin-
independent scattering, and the red region leads to DM overabundance
through freeze-out. See text for further details. 39

3.4 The bounds on the models at two different hierarchies between \( m_\phi \)
and \( m_\chi \). Bounds at 95\% CL from the LHC are obtained from measurements
at 8 TeV and 20 fb\(^{-1}\); the orange (blue) curves depict dilepton bounds
on the RR (RL) models, and are solid (dashed) for \( m_{ee} (\cos \theta_{CS}) \)
bounds; the magenta curves depict jets + MET constraints. The green
curves are 90\% CL Xenon1T constraints on spin-independent scattering,
and the red region leads to DM overabundance through freeze-out.
See text for further details. 40

4.1 Tree level contributions to \( W_\gamma \) production (s, t, and u channels respec-
tively). For leptonically decaying \( W \) one should add an extra diagram
where the photon is emitted by the lepton which corresponds to final
state radiation. 50

4.2 Partonic angular distribution for \( W^\pm \gamma \) production. The RAZ is located
at \( \cos \theta^* = \mp 1/3 \), where \( \theta^* \) is the angle between the incoming quark
and the photon in the CM frame. 51

4.3 Proton level distributions of \( W^+ (\ell^+ \nu) \gamma \) production: Difference of the
photon and lepton rapidities (top), photon rapidity (center), and Collins-
Soper angle (bottom). Solid lines represent the distributions with basic
kinematic cuts (Eqs. 4.2, 4.3), while dashed lines include a cut in the
cluster transverse mass \( m_{T}^{\text{clus}} \) (Eq. 4.1), or the invariant mass \( m_{W_\gamma} \), as
indicated. 53
4.4 Invariant mass distribution of $W^+(\ell^+\nu)\gamma$ production. The blue histogram represents the SM background and colored peaks represent different benchmark signals from a scalar resonance. Solid lines represent full distributions whereas dashed lines are the distributions including a cut in the CS angle for each bin. For the 700 GeV bin, the number of events for both signal and background are explicitly shown by the little numbers next to the bins.

4.5 Production of $W\gamma$ though a charged squirk bound state. A s-up and a s-down quark are produced through a $W$ boson. They confine due to the dark confining force. This bound state is created with an arbitrary configuration of angular momentum $l$. The bound state relaxes very promptly and decays as a scalar resonance.

4.6 Production of $W\gamma$ though scalar (left) and vector (right) resonances.

4.7 Shape of signal distributions for $W^+(\ell^+\nu)\gamma$ resonances. The variable $\Delta y$ has a similar shape for spin 0 and 1 resonances (top). The rapidity $y_\gamma$ tends to be slightly flat for low masses and it tends to focus in the central region for large masses (center). The Collins-Soper angle presents different shapes depending on the spin of the resonance (bottom).

4.8 Ratio variable (left) defined in 4.5 that indicates how much the background decreases by making a cut in the central region of the rapidity distributions. For a given $R_V$, provided the signal shape (Fig. 4.7) this figure shows how much the signal significance (right) Increases.

4.9 Sensitivity of $W\gamma$ scalar resonances for HL-LHC: (Red) Squirk bound states in Folded Supersymmetry for only two generations of squarks, (Blue) Effective Heavy Pions with a confining scale of $\Lambda = 10$ TeV. Solid line provides the bounds using the entire $m_T^{\ell\nu\gamma}$ distribution whereas the dashed line only includes events in the central region of the rapidity distributions.

A.1 Total cross section times branching ratio of the $Z'$ in the di-lepton channel as a function of the $Z'$ boson mass. The interception of the blue, orange and purple lines and the observed limit represent the lower bounds for the $Z'$ mass of our model for the given values of the gauge coupling $g_x$. The black line shows the limit on a Sequential Standard Model (SSM) $Z'$ [131], for comparison.

A.2 Left (Right): Higgs mass as a function of $M_{SUSY}$ ($m_\phi$) for different values of the coupling $g_x$ and fixing $m_\phi = 1$ TeV ($M_{SUSY} = 700$ GeV). The solid lines correspond to the max-mixing case and the dashed lines to the non-mixing case. In both panels we have used the lower bound values for $M_{Z'}$ that correspond to the given $g_x$ values.
A.3 Masses of the stop eigenstates for large $\tan \beta$, max-mixing with a small coupling $g_x = 0.32$ (solid lines) and non-mixing with a larger coupling $g_x = 0.63$ (dashed lines) cases. The gray-shaded region are excluded because $m_{\tilde{t}_2} > m_{\tilde{t}_1}$ and the gray-dotted contours correspond to fixed values of $M_{SUSY}$. The yellow region corresponds to a conservative bound that excludes the lightest stop for $m_{\tilde{t}_1} < 750$ GeV (refer to the text for details).

B.1 $\chi^2$ values from a fit to the $R_{D^{(*)}}$ data, as a function of Wilson coefficients ($C_i$) for the operators generated by LQ exchange.

B.2 Feynman diagrams for Z decay with leptoquark.

B.3 Parameter space of the leptoquark $S_1(\overline{3}, 1, 1)$ in which it can explain the $R_{D^{(*)}}$ anomaly (green shaded region), along with the regions excluded by electroweak precision measurements (above dashed blue line) and pair production bounds (left of solid red line). The solid blue line indicates the electroweak constraints after decoupling the $S_1^{1,2}$. 

References

viii
TABLES

3.1 SIMPLIFIED DARK MATTER MODELS . . . . . . . . . . . . . . . . . . . . . . . 27
B.1 DIMENSION-6 OPERATORS CONTRIBUTING TO $\bar{B} \rightarrow D^{(*)}\tau\bar{\nu}$ . 77
B.2 RELEVANT LEP AND SLC OBSERVABLES . . . . . . . . . . . . . . . . . . . . . . . 85
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Thank you!

(I think I made it, I became a phenomenologist)
CHAPTER 1

INTRODUCTION

This chapter presents a brief retrospect of particle physics to highlight the points in history where different particles were discovered. It also presents the modern view of particle physics describing what is known as the Standard Model (SM). This chapter is intended to introduce the general public to the research developed in this thesis. Here is provided the philosophical approach to the problem of finding new particles and and new forces followed in this thesis in the context of the Large Hadron Collider (LHC) era.

1.1 A Brief Retrospect of Particle Physics

1.1.1 Early particle physics: electrons, protons and neutrons

The first fundamental particle ever discovered was the electron. In 1897 J. J. Thomson determined that cathode rays, i.e. a stream of particles observed in vacuum tubes (those used to build the old box TV), were made by negatively charged particles now called electrons. It was soon understood that electrons are responsible for electric currents inside matter and that their mass is about 1800 times lighter than the lightest atom. The second particle ever discovered was the photon. In 1905 Albert Einstein concluded that light is made out of small packages of energy that behave like particles. These were later called photons and they are understood to be the fundamental constituents of the electromagnetic radiation.

The discovery of radioactivity in 1896 by Henri Becquerel led later to the discovery of new subatomic particles. In a radioactive process, a heavy atom decays into a
lighter one emitting different types of radiation; Gamma, Beta, and Alpha. Gamma radiation is made out of photons, Beta radiation is made out of electrons, and Alpha radiation is made of fully ionized Helium atoms. In 1909 Ernest Rutherford bombarded thin foils of Gold with Alpha radiation. These experiments revealed that atoms have a heavy nucleus with some positively charged and some non-charged particles. The former were called protons. At the beginning, the non-charged particles were hypothesized to be formed by bound states of protons and electrons inside the nucleus but in 1932 Chackwick ruled out such hypothesis discovering the neutron.

1.1.2 The discovery of antimatter

By year 1932 physicists knew that all visible objects in the universe are made out of protons, neutrons, electrons, and photons. Rather than a full picture, that was just the beginning of the story. The phenomenon of atmospheric ionization was attributed to the natural radioactivity of the earth. In 1912 Victor Hess determined that such ionization was caused by radiation of high penetrating power that enters from above into our atmosphere. These are called cosmic rays. Studying cosmic rays, in 1932 Carl Anderson discovered particles with similar characteristics of those of an electron, but with positive charge. This particle, the positron, was the first antimatter particle ever discovered. Antimatter was theorized by Paul Dirac in 1928. Dirac wrote an equation that combines quantum mechanics with Einstein’s special relativity. This equation predicted the existence of positively charged electrons that annihilate with normal electrons to produce electromagnetic radiation. This is the main characteristic of antimatter, and it does not only involve electrons. For example, there can be antiprotons; particles that can annihilate with protons to produce radiation.
1.1.3 Neutrinos and muons

Besides the positron, Anderson and his collaborators also discovered particles with negative electric charge heavier than electrons but lighter than protons. These were later called *muons*. The muons detected by Anderson were produced by highly energetic collisions between cosmic rays and particles in the atmosphere. The muon is a heavier “copy” of the electron, with similar properties (e.g. same electric charge) but with the important difference that the muon is unstable. Muons disintegrate into electrons and some nearly invisible particles called *neutrinos*. Neutrinos were hypothesized by Wolfgang Pauli in 1930 when studying the energy spectrum of the electrons emitted in the process of radioactivity. If a heavy nucleus decays into a lighter one emitting electrons (beta radiation), one can measure the energy of those electrons. By conservation of energy, the beta electrons must carry an energy equal to the difference in mass between the heavy and the lighter nucleus. This comes from Einstein’s equation

\[ E = mc^2, \]

that establishes a relation between energy and mass. The experiments with beta electrons showed a result that suggested the non-conservation of energy and momentum in the process of radioactivity. This led Pauli to formulate the existence of a highly penetrating type of radiation made out of almost non-interacting particles, the neutrino. In Pauli’s idea, neutrinos are produced along with electrons in the beta decay carrying a portion of the energy in the process in a way that energy is conserved. Modern particle physics experiments are able to detect neutrinos from different sources; the atmosphere, the sun, and from deep in the cosmos. The mass of a neutrino has been estimated to be about a million times smaller than the mass of the electron but its exact value has not been determined yet.
1.1.4 The forces of nature

Before 1910, only two forces were known by physicists; gravity and electromagnetism. Forces are tied to an intrinsic property of matter. For example, the force of gravity experienced by an object is related to its mass, whereas the electromagnetic force relates to its electric charge. This means that a neutral and almost massless object like the neutrino cannot feel the electromagnetic force, and it can barely feel the force of gravity. However, the neutrinos can feel another force: the \textit{Weak Nuclear Force} (weak force for short). This force is the one responsible for radioactive decays in which, as previously explained, the neutrino plays an important role.

By 1910, it was evident that besides the force of gravity and the electromagnetism, there must exist another force that binds together the protons inside the atomic nucleus. This force should be strong enough to overcome the electric repulsion of the protons so the nucleus does not blow apart. For this reason, this force is called the \textit{Strong Nuclear Force} (strong force for short). To compare the relative strength of the four forces of nature one can define a scale where the force of gravity between the earth and an average size person is defined as the unit. In such case, the weak force will correspond to $10^{29}$ units (one trillion billion million), the electromagnetic force $10^{40}$ units, and the strong force $10^{43}$ units. Yes, the force of gravity is extremely weak compared to the other three forces of nature.

Forces are understood to be “carried” by mediator particles. In the case of electromagnetism, the photon is the force carrier. As we mentioned, photons can be understood as small packages of energy present in the electromagnetic radiation. Photons play an important role in the description of the interaction between electromagnetic radiation and matter. Because of this, it is possible to detect photons at experiments by impinging gamma radiation in metallic plates and measuring variations in the electric current passing through a circuit attached to the plate. This is called the photoelectric effect. In principle, the force of gravity could be carried by
the graviton, but this particle has never been detected. Given the small strength of gravity, it is difficult to perform experiments to prove the existence of the graviton.

1.1.5 A zoo of new particles

In 1935 Hideki Yukawa predicted the existence of massive particles that mediate the strong nuclear force. In 1947 Cecil Powell et al., discovered Yukawa’s particle using photographic plates to detect cosmic rays. This particle was called the Pion. Even though Yukawa was right about the existence of the pion, he was not totally right for the pions are not the fundamental particles that carry the strong nuclear force. Such job is performed by the gluon, a particle that was discovered in 1978, many years after the discovery of the pion. Another particle detected from cosmic rays was the Kaon in 1947. Something especial about pions and kaons is that they can feel the strong force just like the proton and the neutron.

With the development of the particle accelerators, a plethora of new particles were discovered around the 50’ and 60’. Most of these particles interact or decay through the strong interaction. Few of these new particles were lighter than the proton (these were called mesons) but most of them were heavier (these were called barions). All the new particles had short life times, i.e. they disintegrate very fast into lighter particles. Due to the vast list of new particles, a new organizing principle was necessary. In 1964 Murray Gell-Mann suggested that the strongly interacting particles are composed by quarks; mesons (like pions and kaons) contain two quarks, and barions (like protons and neutrons) contain three quarks. This idea was later adapted by Bjorken and Feynman in 1969, and by the year 1974 the existence of quarks was well accepted in the particle physics community with the discovery of the charm quark (the November Revolution).

\[1\] The actual definition of mesons and barions is provided below.
1.2 Particle Physics Today

1.2.1 Classification of the particles

The Standard Model of Particle Physics suggest that all the visible objects in the universe and most observed phenomena can be understood using the fundamental particles depicted in Fig. 1.1. In the figure, each box shows in the top-left corner the intrinsic properties of each particle (mass, electric charge, and spin). The electric charge is shown in units of the electric charge of the electron e.g. the electron has electric charge of $-1$ (in this case the proton has charge $+1$). The matter particles are divided in *quarks* and *leptons*. The particles that serve as mediators of forces are called *gauge bosons*.

- **Quarks**: These are the fundamental constituents of hadrons. Three quarks form either a proton or a neutron, and two quarks can form either pions or kaons. These particles can feel the strong force. The quarks come in six different flavors (*up*, *down*, *charm*, *strange*, *top*, *bottom*) and are divided in three generations. The first generation contains the *u* (up) and *d* (down) quarks, the second includes the *c* (charm) and *s* (strange), and the third includes the *t* (top), and *b* (bottom) quarks. The up-type quarks (*u, c, t*) differ from the down-type quarks (*d, s, b*) in that the former has electric charge $2/3$ and the latter has charge $-1/3$. When the quarks $uud$ combine, they form a proton with electric charge $+1$.

- **Leptons**: These particles are divided into charged and neutral leptons. Charged leptons include the electron, muon, and the *tau*, a heavier version of the previous two. The neutral leptons are the neutrinos. There is a neutrino for each type of charged lepton (*electron-neutrino*, *muon-neutrino*, and *tau-neutrino*).

- **Gauge bosons**: As we mentioned before the photon and the gluon are the force carriers of the electromagnetic and the strong forces respectively. In the same way, the particles $Z$ and $W$ are the force carriers of the weak force. These particles are massive, as opposed to the photon and the gluon that are massless.

- **The Higgs boson**: This particle appears in the SM as a quanta (or energy package) of what is called the Higgs field. (Similar to the way that photons are the quanta of the electromagnetic field.) The Higgs field is supposed to permeate all the space, similar to the idea of the *ether* in ancient Greece. The Higgs field was formulated as an explanation of the mass of the fundamental particles by Peter Higgs in 1964 (and R. Brout, F. Englert, G. S. Guralnik, C. R. Hagen, T. W. B. Kibble) \[1,2\]. The idea is that when particles interact with
Figure 1.1. The Standard Model of Particle Physics. Each box shows in the top-left corner the mass of each particle, its electric charge, and another property called spin. The mass is provided in units of energy according to Eq. 1.1. The quarks and gluons feel the strong force. Charged leptons, neutrinos, the $Z$ and the $W$ boson feel the weak force. All charged particles can feel the electromagnetic force. The Higgs particle interacts with all massive particles. (Figure from *Wikipedia Commons*.)
the Higgs field, they acquire a mass proportional to the interaction with the field. An analogy of this phenomenon is illustrated in Fig. 1.2. The figure shows different persons trying to swim through a pool. One can imagine that each fundamental particle is represented by a swimmer bearing a different number of floaters. The number of floaters tell us how much interaction there is between a particle (the swimmer) and the Higgs field (the water). The swimmer with a larger number of floaters will move slower in the water. In other words, the particle that interacts the most with the Higgs field will be heavier. The mass of a particle is understood in the subatomic world as a measure of how much such a particle interacts with the Higgs field. The discovery of the Higgs in 2012 by the LHC [4, 5] represents a major achievement of the particle physics community in the understanding of our universe.

1.2.2 Open questions in particle physics

The SM describes most of the phenomena that we observe in the universe. There are however, some questions of the observable universe that the SM can not answer.

- Dark Matter: Astronomical observations have revealed that visible matter (stars, galaxies, etc.) only accounts for about 16% of the matter density of our universe. The other 84% corresponds to matter that we cannot see, i.e. Dark Matter (DM). When astrophysicists measure the velocity of the stars of a galaxy, one expects the velocity of the stars far from the center of the galaxy to be smaller than the velocity of those close to the center. What they observe is that for the farther stars their velocity levels off after certain distance. This
behavior can only be explained if there is additional (invisible) matter in the galaxy.

• Neutrino mass: The SM was originally formulated assuming that neutrinos were massless. Experiments in neutrino physics have revealed that neutrinos must posses a small but non-zero mass. Each neutrino has a different mass, and it is possible that one of them is massless. Particle physicists have not yet been able to measure the exact mass of the neutrinos. For this reason, it is unclear what the mechanism is through which neutrinos acquire mass. The explanation could be as simple as the one provided by the Higgs mechanism, but theoretically, there are more possibilities still open. Some scenarios suggest that not only the Higgs, but new particles could play a role in the mechanism for the neutrino mass. These new particle may have a connection to DM.

• Baryogenesis: It is known that matter and antimatter can annihilate to form radiation, or more accurately, to form gauge bosons (e.g. an electron and a positron can collide to produce two photons or a $Z$). From the theoretical point of view there is no reason why nature would prefer matter over antimatter. This is, at the origin of the universe, when particles were created, it is reasonable to believe that matter and antimatter were produced at the same rate. However, something happened in the early universe that created a matter-antimatter asymmetry, and for this reason the universe we observe today is made out of matter and not antimatter. This process is called Baryogenesis. Physicists have formulated different mechanisms through which Baryogenesis could have taken place. Some of them involve the interaction of the Higgs boson with new particles. Others suggest that DM could have played an important role in the process of Baryogenesis.

• The Flavor puzzle: In the SM an up-type quark can decay into a down-type quark and a $W$ boson, e.g. $t \rightarrow bW^+$. In such a process, the flavor of the quark can change. Because a charged particle is produced (the $W^+$ boson), these processes are called flavor changing charged currents (FCCC). These processes have been detected at particles accelerators and are highly probable. However, processes where a neutral gauge boson is produced and the flavor of the decaying quark is changed, e.g. $t \rightarrow cZ$, do not exist or are highly unlikely in the SM. FCNC is not a problem of the SM but rather a consequence of the flavor structure of the theory. Particle physicists have looked for new interactions and particles that produce flavor changing neutral currents (FCNC) with no success. The flavor puzzle states the following; on one hand, there is evidence that new particles and interactions must exist (DM, neutrino mass, Baryogenesis), but on the other hand, such interactions must respect the flavor structure of the SM and produce undetectable FCNC effects.

The previous examples are probably the most obvious problems in the particle physics community from the phenomenological point of view. There are other problems that
are equally important but from a more theoretical point of view. These include: 1. The *Hierarchy problem* which has to do with the fact that the mass of the Higgs boson is sensitive to the presence of new heavy particles. This can be a problem if the mass of these new particles is too large. This indicates that there might exist new particles that are not much heavier than the Higgs itself and that serve to stabilize the mass of the Higgs, 2. *Unification* is the idea that all forces of nature are different manifestations of a single force that governed the dynamics of the early universe. Particle physicists have been looking for the gauge bosons associated with the *unified force of nature* with no success, 3. The existence of *Extra Dimensions* in the universe still remains a mystery. Einstein’s theory of relativity taught us how time can be understood as a fourth dimension, but we have not found evidence of the existence of extra spatial dimensions.

1.3 The Approach to Particle Physics in this Thesis

The LHC is the largest particle collider ever created. It reaches an energy of 13 TeV which is 13 thousand times higher than the equivalent energy of the mass of a single proton. The most important achievement of the LHC so far has been the discovery of the Higgs boson. The SM predicted the existence of such a particle but it did not predict its mass. The next goal of the LHC is to search for different types of new particles and interactions that appear in the theories formulated to solve some of the problems mentioned above. The LHC searches for different types of exotic new particles: new particles that interact only with matter particles (quarks and leptons), or gauge bosons (photon, $Z$, $W$), or the Higgs, or a combination of these.

In searching for new particles, the basic principle is that the energy of the collision can be used to produce the new particles one is searching for. This is thanks to Einstein’s equation $E=mc^2$. The LHC has reached its highest energy, yet we have not found evidence of any new particle or interaction. This motivates the idea that we
might need to increase the energy of the collider because, maybe, the new particles
we are looking for are too heavy to be produced with the current energy. Projects
for more energetic particle colliders are under consideration and the particle physics
community is moving towards that direction. Waiting while a more powerful collider
is built can take decades. In the mean time, particle physicists have (among others)
one alternative hypothesis to explain why new physics signals have not been detected
yet; it is possible that new particles are actually being produced at the LHC but their
signatures are being overwhelmed by huge backgrounds in our detectors. In other
words, as an expert in telecommunication would say, the signals we are looking for
might be smaller than the noise in the antennas we are using to capture them.

In the search for sneaky signals at the LHC, this thesis uses two approaches:

• High precision measurements: Some processes are expected to be measured with
high precision at the LHC as it accumulates more and more data. An example
is the production of a pair of charged leptons, e.g., the reactions $pp \rightarrow e^-e^+$
or $pp \rightarrow \mu^-\mu^+$. Dilepton signals are clean in the sense that it is relatively
easy to distinguish muons and electrons from other particles at the LHC. This
is because leptons leave an isolated charged track in the detector. Having a
clean signal and a large number of events is what makes dilepton production an
interesting spot to look for new physics; Even the smallest new physics signal
can potentially be detected due to the high precision of these measurements.

• Searches in regions with low noise: Some processes in the SM are predicted to
be either small or nonexistent. One example are the FCNC processes mentioned
above. Another example is found in Diboson production, specifically, in the
reaction $pp \rightarrow W^\pm\gamma$. If one measures the angle between the photon and the
direction in which the center of mass of the $W\gamma$ system is moving, there is an
angular region where little background noise is expected, according to the SM.
Such a kinematic region is a good place to look for new physics signals.

The following chapters show how to unveil new physics using angular distributions
at particle colliders by using these two approaches.
This chapter presents different aspects of particle physics phenomenology and, as such, serves as an introduction to the material that will be discussed in later chapters. In particular, some measurements at the LHC concerning dilepton and diboson production are discussed; an introduction to DM is presented; and the concept of relic density as well as collider signatures of DM are introduced.

2.1 Collider Physics

2.1.1 The Large Hadron Collider

The LHC collides protons at a center of mass energy of 13 TeV and a current integrated luminosity of $\sim 10^2 \text{fb}^{-1}$. It consist of an acceleration system and a series of detectors including the Compact Muon Solenoid (CMS) and A Toroidal LHC Apparatus (ATLAS). At the center of each detector, protons collide in bunches that cross each other every 25 ns, with around 20 collisions per bunch crossing with nominal beam current. This provides of the order of a billion inelastic events per second.

Both detectors, CMS and ATLAS, consist on a series of subsystems, each of which plays a role in identifying the particles that are produced at the primary collisions and in measuring their energy and momentum. These systems are:

- The tracker: This system consist of a series of layers where charged particles can leave their imprints in the form of small electric signals that are amplified and recorded. In the region of the tracker there is a 4 Tesla solenoidal magnetic field with the purpose of bending the path of a particle that passes through
The tracker records the trajectory of the charged particles produced at the collisions and measures their momentum.

- The **electromagnetic calorimeter** (ECAL): This system records the energy of particles that interact electromagnetically and that are light enough to be stopped by the multiple collisions with the atoms of the calorimeter. As a particle travels through the ECAL it scintillates, i.e. produces light, depositing its energy. This light is collected by the electronics and converted into a measurement of the energy of the particle.

- The **hadronic calorimeter** (HCAL): This system is similar to the ECAL, but with a structure of alternating layers of absorber and fluorescent scintillator materials. This makes the HCAL hermetic enough to capture every particle emerging from the collisions (except neutrinos and muons). The energy of the particle entering the HCAL is recorded also from the light collected from the scintillators.

- The **muon system**: Muons can penetrate several meters of iron without interacting. For this reason they are not stopped by the ECAL or the HCAL. The muon system consists of a series of chambers placed at the outermost layer of the detectors. Like the tracker, the muon system uses the powerful solenoidal magnet to measure the energies and momenta of muons produced in primary collisions.

With the use of these systems the LHC can record the identity of particles that are produced at the high energy collisions and measure their momentum and energy. These particles are:

- **Electrons** - leave a signal in the tracker and deposit their energy in the ECAL.
- **Photons** - leave no signal in the tracker and deposit of their energy in the ECAL.
- **Muons** - leave a signal in the tracker and in the muon system.
- **Protons** - leave a signal in the tracker and deposit their energy in the HCAL.
- **Neutrons** - leave no signal in the tracker and deposit their energy in the HCAL.
- **Jets** - strongly interacting particles that deposit their energy in the HCAL. Some of the particles inside a jet leave a signal in the tracker and some others do not.

By measuring the energy and momentum of the particles produced at the primary collisions, it is possible to create a series of kinematic distributions. In the following sections some of these distributions are presented.
2.1.2 The invariant mass distribution in dilepton production

If particles 1 and 2 carrying four-momentum $p_1$ and $p_2$ are produced from a collision, the invariant mass of the process is given by

$$m_{p_1p_2}^2 = (p_1 + p_2)^2.$$ \hspace{1cm} (2.1)

This quantity is equal to the center of mass (CM) energy of the collision ($\hat{s}$). The LHC measures the invariant mass of processes where a pair of oppositely charged leptons of the same flavor ($e^+e^-$ or $\mu^+\mu^-$) are produced. Fig. 2.1 shows some of the recent results of this measurement from the CMS collaboration [6] in the case.
of electron-positron production. Different colors represent the expected number of events from different processes including $t\bar{t}$, $tW$, $VV$ and jets (secondary background) and Drell-Yan $Z/\gamma \rightarrow e^+ e^-$ (signal) as indicated in the figure. The green line is a simulated signal from a model with higher dimensional operators that can produce dileptons.

The invariant mass can be calculated in any process where the four-momentum of the final state particles is measured. The variable $m^2_{ij}$ is not limited to processes with only two final states; for any number $n$ of final state particles the invariant mass has the obvious form $m^2_{p_1...p_n} = (p_1 + ... + p_n)^2$.

2.1.3 The Collins-Soper angle in dilepton production

The differential cross section for $pp \rightarrow l^+ l^-$ calculated in the CM frame depends on the angle between the incoming quark and the outgoing lepton. To calculate the total cross section one assumes that the quark $q$ and the anti-quark $\bar{q}$ carry a given momentum fraction $x_1$ and $x_2$ of the proton momentum. In a hadron-hadron collider it is not possible to know the direction of the incoming quark. For this reason one has to add the probabilities that $q$ carries momentum fraction $x_1$ while $\bar{q}$ carries the fraction $x_2$ with the probability corresponding to the opposite momentum fraction assignation. One then convolutes the partonic differential cross section with the Parton Distribution Functions (PDF) $f_q(x)$:

$$d^2\sigma = \sum_q [f_q(x_1)f_q(x_2) + f_q(x_2)f_q(x_1)]\hat{\sigma}_q dx_1 dx_2,$$

where $f_q(x)$ gives the probability of finding a parton $q$ (quark or gluon) with momentum fraction $x$, and $\hat{\sigma}_q$ is the cross section of the partonic process $q\bar{q} \rightarrow \ell^+ \ell^-$. This procedure makes it hard to keep track of the angular distribution of the final states because of the lack of knowledge of the incoming quark. Nevertheless, it is possible
to define a frame from which one can obtain an angular distribution for the Drell-Yan dilepton production, namely the Collin-Soper (CS) frame [7]. At the LHC most of the proton momentum is carried by the gluons and the quarks. The probability of finding an anti-quark with a momentum fraction higher than the momentum fraction of the quark in a $q\bar{q}$ collision is very small. On this basis one can calculate the cross section separating the two regions where $x_1 < x_2$ and $x_1 > x_2$, and then assigning the larger momentum fraction to the quark.

The LHC measures the CS angle and the forward-backward asymmetry of the dilepton production. The CS angle is defined by

$$\cos \theta_{CS} = \frac{Q_z}{|Q_z|} \frac{2(p_1^+ p_2^- - p_1^- p_2^+)}{|Q| \sqrt{Q^2 + Q_T^2}},$$

where $Q$ is the net momentum of the dilepton system with $Q_z$ ($Q_T$) the longitudinal
(transverse) piece, and $p_t^\pm \equiv (p_t^0 \pm p_t^i) \sqrt{2}$ with $p_1$ ($p_2$) the momentum of the lepton (anti-lepton). Fig. 2.2 shows the results of the CS angle from the ATLAS collaboration for the 8 TeV and 20.3 fb$^{-1}$ data set [8]. The values for which $\cos \theta_{CS} > 0$ ($\cos \theta_{CS} < 0$) represent the forward (backward) region. One can see the forward-backward asymmetry that is present at the parton level in Drell-Yan production. This asymmetry is related to the parity violation in the interactions between the SM fermions and the $Z$ boson.

2.1.4 Transverse momentum and rapidity distributions in diboson production

As mentioned above, the LHC can measure the four-momentum of photons produced from the primary collisions. On the other hand, massive vector bosons are not directly detected. Their presence is inferred from the final states of a given event. For example, if an event is recorded with two oppositely charged leptons and a photon, and the invariant mass of the lepton pair is close to the mass of the $Z$ boson (typically in the range 70-110 GeV), it is inferred that such event corresponds to $Z\gamma$ production. In the same fashion, an event with one lepton, one photon, and missing transverse energy (MET) corresponds to $W\gamma$ production. In this case, it is inferred that the $W$ boson decayed to a charged lepton and a neutrino. The neutrino then escaped invisible to the detector creating an energy imbalance recorded as MET. A third example is $WZ$ production where, for example, one can look for events with three leptons and MET.

The LHC has reported measurements of $WZ$ production [9] in the $\ell^+\nu\ell^-\ell^-$ channel. In this case, the four-momentum of the $Z$ boson, $p_Z^\perp$, is reconstructed from the momenta of the leptons $\ell^+\ell^-$. Among others, this analysis presents two kinematic distributions: the transverse momentum of the $Z$ boson defined as

$$p_T^Z = \sqrt{(p_T^Z)^2 + (p_T^\phi)^2},$$

(2.4)
Figure 2.3. Kinematic distributions for $WZ$ production. The ATLAS collaboration shows the transverse momentum (left) and difference in rapidities of the lepton and the $Z$ boson (right) for 13 TeV and 36.1 fb$^{-1}$. (Plot taken from [9].)

and the rapidity distribution of both the $Z$ boson and the third lepton $\ell'$

$$yz = \frac{1}{2} \log \left( \frac{E_Z + p^Z_T}{E_Z - p^Z_T} \right), \quad y\ell = \frac{1}{2} \log \left( \frac{E\ell + p^{\ell}_T}{E\ell - p^{\ell}_T} \right),$$

(2.5)

where $p^a_i$ corresponds to the component $i$ of the momentum of the particle $a$, and $E_a$ is its energy.

The transverse momentum indicates how much energy is produced from a given collision in the direction perpendicular to the beam axis (defined as the $z$-axis). The collisions at the LHC occur with most of the momentum of the incoming protons along the beam axis. The initial transverse momentum of the partons is too small to cause any important effect. For this reason, the transverse momentum of the final states should add up to zero, according to conservation of momentum.
The rapidities indicate how much the final state particles are boosted along the beam axis. The rapidity of the CM of the final states is called the boost rapidity. A large boost rapidity implies that the momentum fraction carried by one of the partons of a collision is much larger than the momentum fraction of the other parton, \( x_1 \gg x_2 \). According to Eqs. 2.5 if the momentum along the collision axis is much larger than the transverse components \( p_z \gg p_{x,y} \) the rapidity approaches \( \pm \infty \) (because \( E \sim |p| \sim |p_z| \)). Large values of the rapidity are not accessible because of the finite size of the detectors. In the case of photon production, the range of rapidities that can be measured is about \( |y_\gamma| \leq 2.5 \).

In Fig. 2.3 one can see the transverse momentum (left) and the difference in the rapidities of the lepton Z boson and the extra lepton (right). Just like with the dilepton distributions above, the experimental data (black dots) of the diboson distributions agrees with the expectation of the SM (purple histograms).

2.2 Dark Matter

The existence of Dark Matter (DM) is strongly supported by the astrophysical and cosmological data: (i) The velocity distribution of stars within galaxies suggests that these are embedded within a halo of non-visible matter; (ii) A comparison between supernova surveys, baryon acoustic oscillations measurements, and big bang nucleosynthesis with data from the Cosmic Microwave Background reveals that the baryonic matter density represents only one fifth of the total amount of matter in the universe; (iii) The fluctuation spectrum in the distribution of matter in the universe measured by galaxy surveys, lensing surveys, and dust maps, can be explained by assuming the existence of DM. These independent pieces of evidence go across a wide range of energy scales, which makes it hard to discard the existence of DM. The approach to the DM problem in this thesis is the following: If DM exists, how can we detect its signatures at colliders, and can the LHC be the place where the
DM particles are first detected?

2.2.1 The WIMP miracle

By the end of the big bang, the universe is an extremely hot and dense gas. If one assumes that DM has sizable couplings to SM particles and that its number density by the end of the big bang is not zero, then DM gets in thermal equilibrium with the SM particles in the early universe\(^1\). If this is the case, at times when the universe is hot enough, the reaction \(\text{DM} \leftrightarrow \text{SM} \) occurs in both direction so that the number density of DM particles is a constant (up to thermal fluctuations). As the universe expands, it cools down. When the temperature drops below the DM mass, \(T < m_{\text{DM}}\), the SM particles cannot annihilate into DM particles so the number density of the latter undergoes Boltzmann suppression. If at some point, the rate at which the universe expands, \(H(t)\) (the Hubble constant), is faster than the rate at which the reaction \(\text{DM} \rightarrow \text{SM} \) occurs, then DM particles cannot find each other anymore and their number density gets freeze-out (FO)\(^2\).

\(^1\)This assumption is challenged by some models where DM is not a thermal relic, like in the Freeze-In scenario [16].

\(^2\)The actual number density keeps decreasing as the volume of the universe gets bigger and bigger. In reality, the co-moving number density is what gets frozen.
Under reasonable assumptions, the DM relic density can be calculated using \[17\]

\[
\Omega_{\chi}h^2 \simeq 1.07 \times 10^9 \text{GeV}^{-1} \frac{x_f}{M_{\text{Planck}} g^*_1/2 (a + 3b/x_f)},
\]

(2.6)

where \(x_f = m_{\chi}/T_{FO} \sim 25\) is the inverse FO temperature, \(g_* \sim 100\) is the number of relativistic degrees of freedom at FO, \(M_{\text{Planck}}\) is the reduced Planck mass, and the parameters \(a\) and \(b\) can be calculated by expanding the cross section \(\langle \sigma_{\chi}v \rangle \sim a + bv^2\).

WIMP stands for \textit{Weakly Interacting Massive Particles}. The WIMP miracle states that for DM masses close to the electroweak (EW) scale and a coupling between DM and the SM particles of the size of the EW coupling, the formula above predicts a DM relic density close to the experimental value \[18\] \[
\Omega_{\chi}h^2 = 0.1199 \pm 0.0027.
\]

It is a “miracle” because: First, the mass of DM is not known. It could be any number between few \(\text{meV}\) up to the Plank mass (\(10^{19}\) GeV). Second, other evidence for new physics points to the necessity of finding new particles at an energy close to the EW scale (e.g. the hierarchy problem). This means that theories proposed to solve the hierarchy problem can potentially provide DM candidates with EW size masses by invoking the WIMP miracle.

To illustrate the WIMP miracle, one can assume that DM annihilates into neutrinos through \(s\)-channel \(Z\) boson exchange (Fig. 2.4). Then

\[
\langle \sigma_{\chi}v \rangle = \frac{\pi \alpha^2}{2 c_w s_w^4} \frac{1}{m_{\text{DM}}^2},
\]

(2.7)

where the DM-SM coupling is assumed to be the EW coupling \(g\). For DM masses between 0.1 to 1 TeV, Eq. 2.6 reproduces the experimental value for the relic density.
2.2.2 Collider signatures of DM (jets + MET)

Following the idea of the previous example, if DM interacts with the $Z$ boson then it is possible to produce it at colliders. If DM is produced alone, there are no possible signals to keep track of at the LHC because its detectors can not see DM. Instead, one looks for signals where DM is produced along with some other particles, as illustrated in the example in Fig. 2.5. The final states of such a process has a visible particle, the photon, and missing energy. These are called monophoton searches. By measuring the MET distribution of monophoton events the LHC can infer about the existence of DM particles that interact with the $Z$ boson. The lack of a signal does not rule out the existence of DM in general, but it sets constraints on the parameters and assumptions put into the analysis.

In Chapter 3, bounds on DM particles will be set based on searches where two jets are produced in association with large MET. Based on the MET distributions, the LHC sets constraints on DM models that interact with first and second generation quarks in terms of the DM ($\tilde{\chi}^0_1$) mass and the mass of a quark partner (or mediator $\tilde{q}$) that makes possible such interaction [19]. The hypothetical reaction is pair production of the mediators ($pp \rightarrow \tilde{q}\tilde{q}^*$) followed by the decay into DM and jets ($\tilde{q} \rightarrow q\tilde{\chi}^0_1$) as shown in Fig. 2.6. For a given DM and mediator mass, the LHC provides the upper bound on the signal cross section. Any new physics model must produce signals
smaller than what appears in the color code otherwise the LHC would have seen an excess of events. This will be discussed further in chapter 3.

The lack of DM signatures from MET analysis in a variety of processes is one of the motivations for this thesis to focus on DM signatures from quantum interference effects as will be discussed in the next chapter.
CHAPTER 3

DARK MATTER SIGNATURES IN DILEPTON FINAL STATES

This chapter studies a family of simplified models in which DM interacts with both quarks and leptons through renormalizable Yukawa interactions with partner fields. These interactions can reshape the spectra of the lepton pair production at the LHC depending on the properties of the dark sector (e.g. DM spin, mass, couplings). In some regions of parameter space here explored, dilepton distributions place stronger bounds than the ones from other DM probes including direct detection, and jets + MET searches at the LHC. These results were published in [20].

3.1 Motivation

Even though the astrophysical evidence for DM abounds, its fundamental properties remain elusive. Key puzzles that remain unsolved are:

• Is the DM its own antiparticle?
• What is its spin?
• What is its mass?
• How does it couple to the Standard Model, if at all?

These properties result in qualitatively diverse signals in direct detection searches (such as whether scattering is spin-independent, spin-dependent and/or momentum-dependent) and indirect detection (such as whether annihilation is s-wave or p-wave). Thus, different tests may reveal some of the properties of the DM particles [21][27]. This chapter discusses if, under certain assumptions, the same can be done at a
collider. One can find a scarcity in the literature of LHC-related work addressing the questions of self-conjugation, spin and coupling structure, perhaps because the primary focus of most collider searches is to extract the mass of DM and possibly that of a mediator that couples the DM particle to the SM. Some exceptions are Ref. [28], where spin-1/2 and spin-1 DM were distinguished using distributions of MET and jet rapidity; Ref. [29], where DM properties were distinguished by decomposing the missing energy spectrum into basis functions; and Ref. [30], where the matrix element method was employed to distinguish DM spin and mass. See also [31–38]. These previous studies make use of MET, the most striking feature of DM directly produced on-shell at the LHC.

This chapter focuses on collider signals that can potentially address the questions above, but take an approach that is not MET-based. Instead, kinematic distributions with fully visible final states are used. A dark sector can leave its imprint in visible spectra if it induces loop processes interfering with SM amplitudes; in particular, threshold effects may generate distinct signal features. As shown in Ref. [39], such non-resonant signals are best discernible in $\ell^+\ell^-$ production at the LHC: the backgrounds are simple and the rates are high. Indeed, the channel is so clean that in some regions it turns out to be more sensitive to the dark sector parameters than conventional jets + MET and direct detection searches. Note that the collider signals are agnostic to the DM abundance, and can be relevant for models for which the observed particle is only partially responsible for the universe’s DM.

This chapter extends the program of Ref. [39] to study how well dileptonic information can shed light on the quantum properties of the dark sector. This is done using “simplified models” where DM couples to both quarks and leptons, for which appropriately charged mediators are introduced. The simplified frameworks are renormalizable effective theories characterized by a minimal set of inputs, usually no more than the SM-DM coupling, the masses of DM and the mediator, and
specifications of DM spin and the mediator’s quantum numbers [40–54].

While Ref. [39] focused on Dirac DM that only coupled to right handed SM fermions, this chapter surveys and compare several scenarios: DM with spin 0 and spin 1/2, DM that couples to right-handed fermions and left-handed. This analysis of the invariant mass ($m_{\ell\ell}$) and scattering angle data available in dileptonic events. This chapter represents the first study availing LHC measurements of dilepton angular spectra to probe the properties of DM and its mediators. In fact, it is found that angular spectra may provide the strongest constraints in some regions of the parameter space.

3.2 Simplified Dark Matter Models

This study focuses on simplified models in which DM $\chi$ has renormalizable Yukawa interactions with SM fermions $f$ through a partner field $\tilde{F}$, with the interaction schematically given by $L \supset \chi \tilde{F} f$. It is usually assumed that a $Z_2$ symmetry under which all non-SM fields are charged odd (and SM fields charged even) is responsible for DM stability.

Here are considered models comprising two SM singlets $\chi_{A,B}$. A colored field $\tilde{Q}$ is introduced to mediate the singlets’ interactions with quarks, and an uncolored field one $\tilde{L}$ to mediate their interactions with leptons. If $\chi_{A,B}$ are fermions, the mediators $\tilde{Q}$ and $\tilde{L}$ are complex scalars, while if $\chi_{A,B}$ are real scalars the mediators are fermions. Let us consider the following interaction Lagrangian involving these fields:

$$L \supset -\sqrt{2}(\lambda_{\tilde{Q}} \tilde{Q} \chi^+_B q^+_1 + \lambda_{\tilde{L}} \tilde{L} \chi^+_B \ell^+_1) + \text{H.c.},$$

where indices denoting fermion chirality and flavor have been suppressed. For
TABLE 3.1
SIMPLIFIED DARK MATTER MODELS

<table>
<thead>
<tr>
<th>Model</th>
<th>$\chi$ spin</th>
<th>$\tilde{Q}, \tilde{L}$ spin</th>
<th>$\tilde{Q}$ under $G_{SM}$</th>
<th>$\tilde{L}$ under $G_{SM}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$pD^{a}_{RR}$</td>
<td>1/2</td>
<td>0</td>
<td>$(3, 1, 2/3)$</td>
<td>$(1, 1, -1)$</td>
</tr>
<tr>
<td>$pD^{a}_{RL}$</td>
<td>1/2</td>
<td>0</td>
<td>$(3, 1, 2/3)$</td>
<td>$(1, 2, -1/2)$</td>
</tr>
<tr>
<td>$pC^{o}_{RR}$</td>
<td>0</td>
<td>1/2</td>
<td>$(3, 1, 2/3)$</td>
<td>$(1, 1, -1)$</td>
</tr>
<tr>
<td>$pC^{o}_{RL}$</td>
<td>0</td>
<td>1/2</td>
<td>$(3, 1, 2/3)$</td>
<td>$(1, 2, -1/2)$</td>
</tr>
<tr>
<td>$pD^{a}_{RR}$</td>
<td>1/2</td>
<td>0</td>
<td>$(3, 1, -1/3)$</td>
<td>$(1, 1, -1)$</td>
</tr>
<tr>
<td>$pD^{a}_{RL}$</td>
<td>1/2</td>
<td>0</td>
<td>$(3, 1, -1/3)$</td>
<td>$(1, 2, -1/2)$</td>
</tr>
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<td>0</td>
<td>1/2</td>
<td>$(3, 1, -1/3)$</td>
<td>$(1, 1, -1)$</td>
</tr>
</tbody>
</table>

Spin-1/2 DM, the most general DM mass Lagrangian is given by

$$\mathcal{L}_{\text{mass}} = (\chi_A \chi_B) \begin{pmatrix} \delta m & m_\chi \\ m_\chi & \delta m' \end{pmatrix} \begin{pmatrix} \chi_A \\ \chi_B \end{pmatrix} + \text{H.c.} \ . \quad (3.2)$$

A similar-looking (squared) mass matrix may be written down for spin-0 DM in terms of the field $\phi_\chi \equiv (\chi_A + i\chi_B)/\sqrt{2}$ and its conjugate $\phi_\chi^\dagger \equiv (\chi_A - i\chi_B)/\sqrt{2}$.

The fields $\chi_{A,B}$ mix to give mass eigenstates $\chi_{1,2}$, with the lighter species $\chi_1$ serving as DM (referred simply as $\chi$). The mediator masses are free parameters that need not be originated from EWSB. For instance, they may arise from the scalar potential if the mediator is spin-0, or could be vector-like if DM is spin-1/2.

The following assumptions are introduced to simplify the analysis:

1. A common mass $m_\phi \equiv m_\tilde{Q} = m_\tilde{L}$ for the mediators, and equal DM couplings to quarks and leptons, $\lambda \equiv \lambda_{\tilde{Q}} = \lambda_{\tilde{L}}$. 

27
2. DM couples to only a single chirality of SM fermions. This restricts the number of mediator species, since otherwise one would need to introduce mediators that are both singlet and doublet under $SU(2)_W$. Couplings to both left-handed and right-handed leptons are considered, but only couplings to right-handed quarks. Couplings to left-handed quarks are not considered for simplicity. As per fermion generations, couplings to electrons and to either up or down quarks are considered.

3. In order to avoid FCNC, it is assumed the existence of three generations of mediators with their couplings aligned with the SM Yukawa couplings such that, in the mass basis, each mediator generation couples only to a single generation of SM fermions. In order for DM to couple solely to the up/down quark, or to the electrons, it is assumed that mediators of the other generations are heavy.

4. As manifest in Eq. 3.1, it is assumed that only $\chi_B$ interacts with the SM fermions. This assumption captures all the qualitative features of the results here presented; allowing both $\chi_A$ and $\chi_B$ to interact tends to only rescale the couplings required to produce similar signal rates.

5. Setting $\delta m' = 0$ and varying $\delta m$, one can interpolate between Majorana and Dirac (or real and complex scalar) scenarios. Specifically, the Majorana (or real scalar) limit is achieved by tuning $\delta m$, with $\delta m \to \infty$ \[39, 55\], while $\delta m \to 0$ renders spin-1/2 DM Dirac and spin-0 DM a complex scalar. Pure Dirac/complex scalar DM notoriously has a large spin-independent cross section scattering off nuclei, and is excluded by direct detection experiments for the range of DM masses and couplings of interest. Therefore, in this study the limit $\delta m \to 0$ will never be taken, setting $\delta m = 1$ MeV as the lower limit. As discussed in Ref. \[39\], for splittings of this size, DM behaves like a Majorana fermion (if spin-1/2, real scalar if spin-0) in direct detection experiments, since the heavier state is kinematically inaccessible given the local DM velocity $\sim 10^{-3}$. Majorana/real scalar DM typically has a much smaller scattering cross section than the Dirac/complex scalar case and hence is much more viable \[10\] (see Sec. 3.3). Meanwhile, $\mathcal{O}$(MeV) mass splitting is well below the LHC detector resolution, hence $\chi_{1,2}$ are indistinguishable at colliders and DM will appear as a Dirac or complex scalar particle in the collider study. Thus, for a fixed DM mass, varying $\delta m \geq 1$ MeV will have no effect on how DM appears in direct detection as all scenarios will interact as Majorana/real scalars. However, $\delta m$ will dramatically change how DM appears in dilepton distributions. For the remainder of this chapter, the $\delta m \geq 1$ MeV regime will be referred to as “pseudo-Dirac” for spin-1/2 DM and “pseudo-complex-scalar” for spin-0 DM.

6. CP-violating phases in the masses and couplings vanish.

7. Quartic couplings involving new scalars introduced in our set-up are ignored as they have little to no impact on the dilepton signals.

To summarize, it is assumed that DM couples to either right-handed up or down quarks and to electrons of either chirality, with DM itself having spin 0 or 1/2. These
set-ups may be classified into eight models, which are $pD^{u}_{RR}$, $pD^{u}_{RL}$, $pCS^{u}_{RR}$, $pCS^{u}_{RL}$, $pD^{d}_{RR}$, $pD^{d}_{RL}$, $pCS^{d}_{RR}$, and $pCS^{d}_{RL}$. The superscript denotes the quark to which DM couples, and the first (second) subscript the chirality of the quark (lepton), while “$pD$” and “$pCS$” denote whether DM is pseudo-Dirac or a pseudo-complex scalar. The field content of these models is summarized in Table [3.1]

3.3 Dark Matter Signatures at Colliders

In this section are illustrated the various effects of radiative corrections from the dark sector on $pp \rightarrow \ell^+\ell^-$ spectra, and how these may help distinguish the properties of $\chi$. One can go about this task by contrasting the signals produced by mutually exclusive cases of a single property, e.g. comparing signals of $pD^{u}_{RR}$ and $pCS^{u}_{RR}$ while keeping all masses and self-conjugation properties the same.

Assuming massless quarks and leptons, and denoting by $\theta$ the center-of-momentum scattering angle between the incoming quark and outgoing lepton, the parton level leading order (LO) Drell-Yan double differential cross section is given by

\[
d\sigma_{\text{tot}} = \frac{d^2\sigma_{\text{tot}}}{dc_{\theta} dm_{\ell\ell}} = d\sigma_{\text{SM}} + d\sigma_{\text{int}} + d\sigma_{\chi}, \tag{3.3}
\]

with

\[
d\sigma_{\text{SM}} = \frac{1}{32\pi m_{\ell\ell}^2 N_c} \sum_{\text{spins}} |\mathcal{M}_{\text{SM}}|^2, \\
d\sigma_{\text{int}} = \frac{1}{32\pi m_{\ell\ell}^2 N_c} \sum_{\text{spins}} 2\text{Re}(\mathcal{M}_{\text{SM}}^* \mathcal{M}_{\chi}), \\
d\sigma_{\chi} = \frac{1}{32\pi m_{\ell\ell}^2 N_c} \sum_{\text{spins}} |\mathcal{M}_{\chi}|^2 \tag{3.4}
\]

where $N_c = 3$ is the number of QCD colors, $\mathcal{M}_{\text{SM}} = \mathcal{M}_{\gamma} + \mathcal{M}_{Z}$ is the SM amplitude for the tree-level Feynman diagram in Fig. [3.1]
Figure 3.1. Feynman diagrams for dilepton production at the LHC. On top is the Standard Model Drell-Yan process at tree level. The middle row shows the box contributions from pseudo-Dirac DM with scalar mediators. The bottom row shows the same from pseudo-complex DM with fermion mediators. See the text and Table 3.1 for more details.

As the NP effects enter at loop level, care must be taken to ensure that all effects at a given coupling order are consistently included. Additionally, purely SM loop (mainly QCD) effects must be accounted for. These issues give rise to the following considerations:

- One-loop SM effects enter at the amplitude level at $O(g^2 g_s^2)$, where $g$ and $g_s$ are the QED and QCD couplings, while NP effects enter at $O(g^2 \lambda^2)$ for vertex corrections and $O(\lambda^4)$ for the box diagrams. The net result of the purely SM loop effects is to replace $d\sigma_{\text{SM}}$ in Eq. (3.4) by the SM cross section at next to leading order (in QCD), $d\sigma_{\text{SM,NLO}}$.

- Interference between SM and NP loops ($d\sigma_{\text{int}}$) results in contributions to $d\sigma_{\text{tot}}$ of $O(g^4 \lambda^2)$ and $O(g^2 \lambda^4)$, where the former involve vertex corrections and the latter involve box diagrams. Comparing these terms, one finds the box contributions significantly larger when $\lambda \sim 1$, which is also the regime of couplings where the NP effects have enough statistical significance for LHC bounds to apply.
This happens due to the difference in power-counting the couplings, and also because the box diagrams generate more pronounced threshold effects.

- The $d\sigma_\chi$ term involves the square of box and vertex corrections. These are, in principle, the same order in perturbation theory as the interference between the tree-level and NP two-loop amplitudes. As NP effects have been calculated only at one loop, most terms in $d\sigma_\chi$ cannot be consistently included in the calculation. An important exception that can be included consistently is the square of the NP box diagrams, which is the only $O(\lambda^8)$ contribution to the cross section at any order.

As the box diagrams dominate the interference term, in the following discussion the vertex correction will be dropped and use “$\mathcal{M}_\chi$” as a loose notation to describe the box amplitude. In order to provide the resulting cross section expressions $d\sigma_{\text{int}}$ and $d\sigma_\chi$ it is convenient to define the following short-hand notation for the Passarino-Veltman (PV) box functions:

$$D_i \equiv D_i[m_{q_1}^2, m_{q_2}^2, m_{l_1}^2, m_{l_2}^2, \hat{s}, \hat{t}, \mu_1^2, m_{\phi_1}^2, m_{\phi_2}^2], \quad (3.5)$$

where $i$ is the PV index, and $\mu_{1,2}$ are the DM eigenmasses of Eqs. 3.2 and ???. Since in Sec. 3.2 it has been set $\mu_2 - \mu_1 = 1$ MeV, which is unresolvable at the LHC, one may well approximate $\mu_1 = \mu_2 = m_\chi$. In general, the interference between the tree-level and box amplitudes can be split into a piece in which the tree-level diagram has an $s$-channel–mediated photon, and another in which it has an $s$-channel–mediated $Z$ boson: $d\sigma_{\text{int}} = d\sigma_{\gamma-\text{box}} + d\sigma_{Z-\text{box}}$. And mentioned, one can approximate $d\sigma_\chi$ with the cross-section coming from the squared box amplitude, $d\sigma_{\text{box-\text{box}}}$. In the following expressions for these cross sections for our various models are provided, up to a proportionality factor $(32\pi m_{\ell_1}^2 N_c)^{-1}$. Here $e$ is the QED coupling, $Q_q = 2/3(-1/3)$ is the electric charge of up-type (down-type) quarks, $g$ and $c_W$ are the electroweak coupling and mixing angle respectively, $a_f$ and $b_f$ are respectively the vectorial and

\footnote{E.g. the cross term between the NP vertex correction and the box diagrams is $O(g^2\lambda^6)$, the same as the interference between the tree-level SM and a NP two loop amplitude.}

31
axial couplings between SM fermions and the $Z$ boson, and $m_Z$ and $\Gamma_Z$ are the mass and the width of the $Z$ boson. These expressions were obtained using FeynCalc [56, 57], and numerical results obtained with LoopTools [58] and Package-X [59].

3.3.1 Cross sections for RR models

The signal cross sections for the $pD^u_{RR}$, $pD^d_{RR}$, $pCS^u_{RR}$, $pCS^d_{RR}$ models are given by

$$
d\sigma_{\gamma-box} \propto -\frac{e^2 Q_q (\hat{s} + \hat{t})^2 \lambda^4}{4\pi^2 \hat{s}} \text{Re}[\tilde{D}],
$$

$$
d\sigma_{Z-box} \propto -\frac{g^2 (a_\ell - b_\ell)(a_q - b_q) (\hat{s} + \hat{t})^2 \lambda^4}{16\pi^2 c_W^2 (m_Z^2 - \hat{s})^2 + m_Z^2 \Gamma_Z^2} \times \text{Re} \left[ (m_Z^2 - \hat{s} - im_Z\Gamma_Z) \tilde{D} \right],
$$

$$
d\sigma_{\text{box-box}} \propto \frac{(\hat{s} + \hat{t})^2 \lambda^8}{64\pi^4} |\tilde{D}|^2,
$$

(3.6)

where $\tilde{D}$ are combinations of PV functions:

$$
\tilde{D}_{pD_{RR}^{u,d}} = 2D_{00} + \hat{s} (D_2 + D_{12} + D_{22} + D_{23}),
$$

$$
\tilde{D}_{pCS_{RR}^{u,d}} = 2D_{00} - \hat{t}D_{13}.
$$

(3.7)

3.3.2 Cross sections for RL models

The signal cross sections for the $pD^u_{RL}$, $pD^d_{RL}$, $pCS^u_{RL}$, $pCS^d_{RL}$ models are given by

$$
d\sigma_{\gamma-box} \propto -\frac{e^2 Q_q t^2 \lambda^4}{4\pi^2 \hat{s}} \text{Re}[\tilde{D}],
$$

$$
d\sigma_{Z-box} \propto -\frac{g^2 (a_\ell + b_\ell)(a_q - b_q)t^2 \lambda^4}{16\pi^2 c_W^2 (m_Z^2 - \hat{s})^2 + m_Z^2 \Gamma_Z^2} \times \text{Re} \left[ (m_Z^2 - \hat{s} - im_Z\Gamma_Z) \tilde{D} \right],
$$

$$
d\sigma_{\text{box-box}} \propto \frac{t^2 \lambda^8}{64\pi^4} |\tilde{D}|^2,
$$

(3.8)
where
\[
\tilde{D}_{p_{RL}^{\alpha \delta}} = m_{\chi}^2 D_0, \\
\tilde{D}_{p_{CSR}^{\alpha \delta}} = 2 D_{00} - (\hat{s} + \hat{t}) D_{13}.
\]

As the focus in this section is on the qualitative differences between various DM models, rather than between the SM and DM, the following results will be calculated using \(d\sigma_{SM,LO}\) for now. The inclusion of NLO SM effects and the considerations here itemized will be used in Sec. 3.2 when the dilepton distributions are used to derive limits on the DM parameter space.

3.3.3 Qualitative results

The most unique feature of \(d\sigma_{tot}\) occurs at \(\sqrt{s} \gtrsim 2m_{\chi}\), when \(\chi\) goes on-shell in the box diagrams in Fig. 3.1 and \(M_{\chi}\) develops an imaginary part \(\text{Im}(M_{\chi})\) determined by the optical theorem. The \(\text{Im}(M_{\chi})\) is proportional to the product of the amplitudes of the tree-level diagrams (with \(\chi\)'s and fermions as external legs) obtained from “cutting” the box diagram vertically. This imaginary part feeds into the real part \(\text{Re}(M_{\chi})\) through dispersion relations, causing the amplitude to rapidly rise near the threshold. At \(\sqrt{s} \gg 2m_{\chi}\), \(\text{Re}(M_{\chi})\) falls away while \(\text{Im}(M_{\chi})\) takes over as the dominant contributor to \(|M_{\chi}|^2\). The net effect of this takeover at \(\sqrt{s} \gg 2m_{\chi}\) is no more than the addition of a new channel of dilepton production, hence \(d\sigma_{tot}\) will be separated from \(d\sigma_{SM}\) by some offset. All these effects are reviewed in detail in Ref. [39], where the shape of the new physics spectrum was identified as a “monocline.” (See also Ref. [60], which comprehensively reviews dispersion relations.) In the following, it is shown that the above effects also carry the imprint of DM’s microscopic properties, leading to diverse features in dilepton spectra.

To obtain the \(m_{\ell\ell}\) spectra, one integrates the cross sections in Eq. 3.3 over \(\cos \theta\),
Figure 3.2. Dilepton invariant mass (left-top), CS angle (right-top), and angular asymmetry distributions (bottom). The blue line represents the SM background from Drell-Yan production. The orange, green, red, and purple represent the pseudo-Dirac, pseudo-complex, Majorana, and real scalar cases. Solid lines represent models with right-handed quarks and right-handed leptons (RR) whereas dashed lines represent the RL models. Signal lines are plotted at the benchmark point $\lambda = 2.0$, $m_\chi = 500$ GeV, and $m_\phi = 550$ GeV.

and for the CS angle one integrates them over $400$ GeV $\leq m_{\ell\ell} \leq 4500$ GeV, the range used by the 8 TeV ATLAS analysis [61]. It is often useful to characterize the angular spectrum as a forward-backward asymmetry,

$$A_{FB} \equiv \frac{N(c_\theta > 0) - N(c_\theta < 0)}{N(c_\theta > 0) + N(c_\theta < 0)},$$  \hspace{1cm} (3.9)$$

or a center-edge asymmetry,

$$A_{CE} \equiv \frac{N(|c_\theta| < \cos \theta_0) - N(|c_\theta| > \cos \theta_0)}{N(|c_\theta| < \cos \theta_0) + N(|c_\theta| > \cos \theta_0)},$$  \hspace{1cm} (3.10)$$
which marks out how much scattering occurs in central regions.

The above discussions are now applied to the various models previously introduced. All spectra are shown by convolving parton-level cross sections with MSTW2008NLO parton distribution functions (PDFs) \cite{62} at $\sqrt{s} = 13$ TeV. For the illustrative plots here, one can approximate LHC dilepton production with the Drell-Yan process $q\bar{q} \rightarrow \ell^+\ell^-$. The treatment of secondary processes that also contribute to dilepton production, such as diboson, $t\bar{t}$, dijet and $W$+jets, will become important when setting constraints in Sec. 3.4.

An illustrative benchmark point with the coupling $\lambda$ fixed to 2.0, and masses $m_\chi = 500$ GeV and $m_\phi = 550$ GeV is chosen for the analysis in this section. As elaborated in Ref. \cite{39}, varying $\lambda$ has the effect of raising or lowering $d\sigma_{\text{int}}$ and $d\sigma_\chi$. Depending on the sign of $d\sigma_{\text{int}}$, this could enhance or diminish the DM signal. Moreover, increasing (decreasing) $m_\phi$ enhances (diminishes) the bump feature near $m_{\ell\ell} \simeq 2m_\chi$. These variations affect the signal significance at the LHC, a point that will be discussed when finding the constraints in Sec. 3.4. Note that the spectrum chosen here, being a “compressed” one, is illustrative of a point where the dilepton probes are expected to outperform jets + MET searches, which suffer from low signal acceptance in these regions.

Finally, in computing the dilepton distributions the following kinematic cuts are being used:

$$|\eta_\ell^\pm| \leq 2.4, \quad p_T^{\ell^\pm} \geq 40 \text{ GeV}.$$  

(3.11)

In the following sections a sketch and contrast of the spectral features induced by various DM species is presented. It is also shown why differences arise between mutually exclusive cases (e.g. spin-0 vs spin-1/2 DM), and explain how these differences can help to sort out the properties of DM and the mediators.
3.3.4 Self-conjugation

As explained in Sec. 3.2, one may interpolate between the Dirac (complex) and Majorana (real) limits of DM by tuning $\delta m$. These limits are readily distinguished by the monocline signature, as shown in Fig. 3.2, where it is plotted $d\sigma_{\text{tot}}$ in the non-self-conjugate limit $\delta m \to 0$ for the models $pD_{u_{\text{RR}}}$ (solid orange), $pD_{u_{\text{RL}}}$ (dashed orange), $pCS_{u_{\text{RR}}}$ (solid green) and $pCS_{u_{\text{RL}}}$ (dashed green), as well as at the self-conjugate limit $\delta m \to \infty$ for the models $pD_{u_{\text{RR}}}$ (solid red) and $pCS_{u_{\text{RR}}}$ (solid purple). In the self-conjugate limit, a subdued signal is produced, while non-self-conjugate DM can produce large, detectable signals. Also, in the self-conjugate limit $pD_{u_{\text{RR}}}$ ($pCS_{u_{\text{RL}}}$) gives near-identical cross sections as $pD_{u_{\text{RR}}}$ ($pCS_{u_{\text{RL}}}$). One may compare all these signals with the blue curve, which corresponds to $d\sigma_{\text{SM}}$.

In the limit $m_\phi, m_\chi \gg \hat{s}$, where the loops can be shrunk to contact operators, the suppressions in the self-conjugate limit are consistent with the loop functions given in the effective theory treatment of Ref. [63]. Since the suppressed rates are a result of a modest addition to $\mathcal{M}_{\text{SM}}$ from the dark sector in the self-conjugate limit, no sizable signals appear in the angular spectra either. Finally, here it is re-emphasized that the “non-self-conjugate limit” does not correspond to Dirac or complex scalar DM, but only to the limit where $\delta m$ is small enough to be irresolvable at colliders while remaining large enough to evade direct detection constraints$^2$.

3.3.5 Spin

In Fig. 3.2, one can see a pronounced “kick” in the signal rates at $m_{\ell\ell} \simeq 2m_\chi$ for fermionic DM ($pD_{u_{\text{RR}}}, pD_{u_{\text{RL}}}$), while the rise in rates appears gentle for scalar DM $^2$

---

$^2$Of course, if the stabilizing $Z_2$ symmetry were broken such that $\chi_1$ decays well within the lifetime of the universe, a pure Dirac or complex scalar formed with $\chi_{1,2}$ is viable. Then $\chi_1$ is no longer the galactic dark matter searched for at direct detection, and only collider constraints apply. The decay length of $\chi_1$ determines whether MET + X or a displaced vertex is the relevant signature. In all cases our dilepton signatures apply, though the effect of non-trivial widths must now be carefully treated.
$(p_{\text{CS}_{\text{RR}}}^\nu, p_{\text{CS}_{\text{RL}}}^\nu)$. This may be understood from the fact that near threshold, the box amplitude is determined by $\text{Im}(\mathcal{M}_\chi)$, which, as mentioned above, is in turn determined by the tree-level amplitudes for $f\bar{f} \to \chi\chi$. The pair-production of complex scalar $\chi$ is more phase-space suppressed at threshold than a Dirac $\chi$, ultimately resulting in a subdued slope of the rise in $d\sigma/dm_{ee}$ for spin-0 DM. In any case, even this difference fades for larger mass splittings between the mediators and DM, where the kick feature is not as pronounced.

One would naively expect the angular distributions to discern the spin of DM, on the strength of their ability to clearly distinguish the spin of mediators in the $s$-channel [64] and $t$-channel [65]. However, while angular spectra are capable of picking up the spin of new particles interfering with the SM via tree-level amplitudes, the angular spectrum resulting from interference with a loop amplitude is non-trivial. As the loop consists of particles with multiple spins, the information on the spins is washed away in the spectrum.

3.3.6 Mass

From the previous sub-section, it is apparent that the mass of DM may be readily cornered if DM is a fermion and if its mass is not much separated from the mediator’s. In that case, the pronounced kick feature in the $m_{\ell\ell}$ signal appears at an invariant mass of $2m_\chi$. As this feature is a result of amplitude-level deviations, it must also be reflected in some way in angular observables plotted as a function of $m_{\ell\ell}$. For instance, one would see it in the $A_{\text{FB}}$, defined in Eq. (3.9), plotted at the parton level (for illustration) in the left-bottom panel of Fig. 3.2 using the same color code as above. The behavior of the $A_{\text{FB}}$ as a function of $m_{\ell\ell}$ with respect to the SM is in accord with the behavior of the $m_{\ell\ell}$ spectrum – the telltale imprint of interference effects. Consequently, an abrupt change of slope is visible in the orange curves at $m_{\ell\ell} \simeq 2m_\chi$. One would also see the kick feature in $A_{\text{CE}}$ (defined in Eq. 3.10)
plotted at the partonic level in the right-bottom panel of Fig. 3.2 where the choice \( \cos \theta_0 = 0.596 \) sets the SM value to zero. Once again the abrupt change of slope at \( m_{\ell\ell} \approx 2m_\chi \) may be seen in the orange curves.

3.3.7 Chirality

The relative chirality between the quarks and leptons in the new physics amplitude, i.e., whether the model is \( \text{RR} \) or \( \text{RL} \), shows up in dilepton spectra in quite interesting ways. (\( \text{RR} \) and \( \text{LL} \), and separately \( \text{RL} \) and \( \text{LR} \), yield similar spectra.) It is known from \( Z \) boson physics and from contact operator analyses (such as in [61]) that the chiral nature of new states has an impact on the scattering angle. Such an impact is seen in the cases where the NP amplitude interferes constructively with the tree-level one, i.e. in \( pCS_{\text{RL}}^u \) and \( pCS_{\text{RR}}^u \). In Fig. 3.2 more forward (\( \cos \theta_{CS} > 0 \)) events are produced by \( pCS_{\text{RR}}^u \) versus \( pCS_{\text{RL}}^u \), and \( pCS_{\text{RL}}^u \) produces visibly more backward (\( \cos \theta_{CS} < 0 \)) events than \( pCS_{\text{RR}}^u \). These differences are best seen by plotting the \( A_{FB} \), as in the left-bottom panel of Fig. 3.2. The \( A_{FB} \) neatly separates the cases of \( \text{RR} \) and \( \text{RL} \), putting them above and below the SM value at large \( m_{\ell\ell} \). Further, the \( A_{CE} \) in the right-bottom panel of Fig. 3.2 clearly signals the chirality combination by putting \( \text{RR} \) (RL) above (below) the SM value.

3.4 Constraints on the DM Parameter Space

Having illustrated that dilepton distributions may help distinguish the properties of DM, one can see how DM could in fact reveal itself first in LHC measurements of dilepton events. This section presents the constraints on the models introduced above from the available LHC data on \( m_{ee} \) and \( \cos \theta_{CS} \). These constraints are compared with those from conventional DM searches such as jets + MET, direct detection, and relic density measurements.

The limits are shown in the plane of the Yukawa coupling \( \lambda \) versus DM mass \( m_\chi \).
Figure 3.3. The bounds on the models at two different hierarchies between \( m_\phi \) and \( m_\chi \). Bounds at 95% CL from the LHC are obtained from measurements at 8 TeV and 20 fb\(^{-1}\); the orange (blue) curves depict dilepton bounds on the RR (RL) models, and are solid (dashed) for \( m_{ee} \) (cos \( \theta_{CS} \)) bounds; the purple regions are excluded by jets + MET searches. The green curves are 90% CL Xenon1T constraints on spin-independent scattering, and the red region leads to DM overabundance through freeze-out. See text for further details.

in Figs 3.3 and 3.4 which correspond respectively to spin-1/2 and spin-0 DM; the left-hand (right-hand) panels depict DM coupling to up (down) quarks. As mentioned in the previous section, these limits must depend on the hierarchy of mediator and
Figure 3.4. The bounds on the models at two different hierarchies between $m_{\phi}$ and $m_{\chi}$. Bounds at 95% CL from the LHC are obtained from measurements at 8 TeV and 20 fb$^{-1}$; the orange (blue) curves depict dilepton bounds on the RR (RL) models, and are solid (dashed) for $m_{ee}$ ($\cos \theta_{CS}$) bounds; the magenta curves depict jets + MET constraints. The green curves are 90% CL Xenon1T constraints on spin-independent scattering, and the red region leads to DM overabundance through freeze-out. See text for further details.

DM masses. Thus we pick two benchmark spectra for illustrating the constraints, one where the spectrum is “compressed” with $m_{\phi} = 1.1 \ m_{\chi}$, and one where it is “uncompressed” with $m_{\phi} = 2 \ m_{\chi}$. These correspond to the top and bottom row
respectively. Throughout the analysis $\delta m = 1$ MeV is fixed. As explained in Sec. 3.2, this gives DM safety from direct detection constraints while masquerading as a Dirac/complex scalar particle at the LHC.

3.4.1 Dilepton constraints

The orange and blue curves in Figs 3.3 and 3.4 show the 95% CL limits from the LHC Run 1 (8 TeV, 20 fb$^{-1}$) measurements of dilepton spectra, corresponding to the RR and RL respectively. The solid (dashed) curves correspond to $m_{ee} \cos \theta_{CS}$ measurements by ATLAS [61, 66] in $e^+e^-$ production. Due to the similarity of results, one can expect similar limits from CMS data [67, 68] and from dimuon production. A $\Delta \chi^2$ fit is performed as done in [39], but considerably improve on the treatment to obtain realistic bounds.

Broadly speaking, the recasting of dilepton measurements into bounds is performed by comparing between three sets of events across $m_{\ell\ell}$ or $\cos \theta_{CS}$ bins (labelled by $i$): the data $N_{d_i}$, the background $N_{b_i}$, and the signal $N_{s_i}$. We take $N_{d_i}$ from ATLAS [61, 66]. It is useful to divide the background into its dominant and subdominant components, $N_{b_i} = N_{b_i}^{\text{dom}} + N_{b_i}^{\text{sub}}$. The former comprises of the Drell-Yan $s$-channel process in Fig. 3.1, while the latter (which we also take from ATLAS) comprises of the reducible backgrounds of diboson, top, dijet, and $W + \text{jets}$ production.

To obtain an $N_{b_i}^{\text{dom}}$ that is as accurate as possible, some kinematic cuts has been imposed as described in Sec 3.3, and Drell-Yan events are obtained at NLO-QCD (i.e. at $O(\alpha_s)$) using MCFM8.0 [69] with MSTW2008NLO PDFs [70], and a renormalization and factorization scale of $m_{\ell\ell}$. Then we account for the efficiency of lepton reconstruction by scaling our events by a global factor that best matches the Drell-Yan background provided by ATLAS. At $\sqrt{s} = 8$ TeV, this factor is 0.74 (0.67) for the $m_{ee} \cos \theta_{CS}$ distribution.

Obtaining the signal events is a subtler process. First, one should obtain the par-
ton level total cross section $d\sigma_{\text{tot}}^i$ defined in Eq. (3.3). For reasons explained in Sec. 3.3, one can neglect two terms: the interference between the SM $\mathcal{O}(\alpha_s)$ corrections and $\mathcal{M}_X$, and all terms involving triangle diagrams. Next, one can convolve $d\sigma_{\text{SM}}^i$ and $d\sigma_{\text{tot}}^i$ with MSTW2008NLO PDFs to obtain the hadron-level cross sections $d\tilde{\sigma}_{\text{SM}}^i$ and $d\tilde{\sigma}_{\text{tot}}^i$. The $N_{s_i}$ are now obtained by first scaling the dominant background by $d\tilde{\sigma}_{\text{tot}}^i/d\tilde{\sigma}_{\text{SM}}^i$, and then adding the result to the subdominant background:

$$N_{s_i} = N_{b_i}^{\text{dom}} \left( \frac{d\tilde{\sigma}_{\text{tot}}^i}{d\tilde{\sigma}_{\text{SM}}^i} \right) + N_{b_i}^{\text{sub}}.$$  

Using all the above information, one can compute

$$\chi_s^2 = \sum_{i=1}^{N_{\text{bins}}} \frac{(N_{d_i} - N_{s_i})^2}{N_{s_i} + \delta_{\text{sys}_i}^2},$$

$$\chi_b^2 = \sum_{i=1}^{N_{\text{bins}}} \frac{(N_{d_i} - N_{b_i})^2}{N_{b_i} + \delta_{\text{sys}_i}^2},$$

and locate the 95% CL bound at $\Delta \chi^2 = \chi_s^2 - \chi_b^2 = 5.99$. Here the systematic errors $\delta_{\text{sys}_i}$ are taken from [61, 66].

The central findings are best understood by directly comparing the right- and left-hand panels of Figs 3.3 and 3.4. The relative behavior of these bounds is dictated by two ingredients – (i) the PDFs: as the up quark has higher parton densities in the proton than the down quark, one expects stronger dilepton bounds for DM coupling to up quarks for DM coupling to down quarks, and (ii) interference effects, or more precisely, the signal contribution of the interference versus the squared box, i.e. $d\sigma_{\text{int}}$ versus $d\sigma_X$ in Eq. (3.4). For example, for the models $pD_{\text{RR}}^u$ and $pD_{\text{RL}}^u$, the tree-level and box diagrams interfere destructively, resulting in a deficit of events with respect to the SM for $m_{\ell\ell} < 2m_X$; this may be seen in Fig 3.2. (On the other hand, the relative sign of the down quark’s electric charge with respect to the up quark ensures that tree-box interference in the case of DM coupling to down quarks is constructive.)
As this interference effect occurs in the low $m_{ee}$ bins, where the event population is high, its contribution to the signal $\chi^2$ could be considerable. This explains why the $m_{ee}$ bound for $pD^u_{RR}$ (and to some extent $pD^u_{RL}$) is so much stronger on the left-hand than on the right-hand panels of Fig. 3.3. At the same time, the $\cos \theta_{CS}$ bounds do not show this hierarchy since the effects of the deficit below and excess above $m_{ee} \simeq 2m_\chi$ are washed out by the integration over $m_{ee}$ bins.

In all four plots, one finds the RR models more constrained than the RL models. In the spin-1/2 DM models this can be understood from the observation made in Sec. 3.3, that $pD^u_{RL}$ gives smaller cross sections than $pD^u_{RR}$ due to differences in how interference proceeds between the standard and crossed boxes. As for the spin-0 DM models, one sees from Fig. 3.2 that $pCS^u_{RR}$ yields slightly larger cross sections than $pCS^u_{RL}$ and is thus subject to slightly stronger constraints. Also, as discussed in Sec. 3.3, the dilepton signal rates decline with $m_\phi/m_\chi$ due to propagator suppression in the loop. This results in the weaker limits in the $m_\phi = 2m_\chi$ plots in comparison to the $m_\phi = 1.1 m_\chi$ plots: in fact, the $pD^u_{RL}$ limits are so weak as to disappear from the parametric range displayed.

3.4.2 Other DM probes

One can now compare the dilepton results with conventional DM probes. In addition to modifying dilepton spectra, the models here studied also have the following:

(a) They can pair-produce colored mediators both through QCD and through exchanging $\chi$ in the $t$-channel of a $q\bar{q}$-initiated process, and these mediators can decay to a quark and DM. Thus, these models confront constraints from dedicated searches for the mediators using jets plus missing energy signatures,

(b) DM can annihilate into quarks and leptons through $t$-channel exchange of mediators and freeze out in the early universe, confronting the relic density measurement by Planck.

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3Ref. 39 had incorrectly flipped the sign of $d\sigma_{int}$ for the model $pD^u_{RR}$, and had derived a bound weaker than that in this work.
(c) DM can scatter against nucleons through $s$-channel exchange of the mediator $\tilde{Q}$, confronting underground direct detection searches.

Bounds derived as a result of (a) - (c) can be seen in Figs 3.3 and 3.4 along with the dilepton bounds explained earlier.

Turning first to the jets + MET bounds, the purple regions are excluded at 95% CL by the CMS Run 1 search [19]. To determine this bound, one can reinterpret the T2qq bounds that assume squark production through QCD (the gluino is decoupled) followed by prompt decay to light quark + LSP. The leading-order cross sections for $\tilde{Q}$ pair production were obtained using MadGraph5 [71] and CTEQ6L1 PDFs [72], and matched with exclusion cross sections provided by CMS. An assumption here made was that detector acceptances of the DM models are similar to the T2qq model of CMS. Since this constraint is agnostic to the chirality of the lepton in our models (with both spin-0 and spin-1/2 DM), we do not distinguish between RR and RL.

The red curve in the plots corresponds to the thermal line where $\Omega_\chi h^2 = 0.12$, with DM being overproduced in the red shaded region below. This curve was obtained using Micromegas4.3 [73], and takes into account co-annihilation between DM and mediators, which becomes important in the compressed region $m_\phi \lesssim 1.1 m_\chi$ [44].

Leptonic modes constitute only a small fraction of the annihilation cross section $\langle \sigma v \rangle$, as opposed to quark modes that come with a color factor of 3. Therefore, though the RL models must give a slightly higher $\langle \sigma v \rangle$ than the RR models due to neutrino final states, there is no visible difference in the red curves. Note that the DM in the models here presented make up a fraction of the total DM population in regions where most of our dileptonic bounds apply.

Finally, the green curves show 90% CL bounds from spin-independent scattering at Xenon1T [74]. (The current spin-dependent limits are consistently weaker and are

4Co-annihilation also occurs between the eigenstates $\chi_1$ and $\chi_2$, but as their mass splitting $\delta m = 1$ MeV $\ll m_\chi$, this practically amounts to the self-annihilation of Dirac/complex scalar DM.
To obtain these bounds, one can assume that the density fraction of DM at freeze-out equals the density fraction in the galactic halo today, i.e. \( \Omega \chi h^2/(0.12) = \rho \chi/(0.3 \text{ GeV cm}^{-3}) \), which effectively scales the exclusion cross sections by \( 0.12/\Omega \chi h^2 \).

The annihilation and scattering cross sections are provided below, Eqs. 3.13 and 3.16.

The dileptonic probes are highly complementary to jets + MET searches. At \( m_\phi = 1.1m_\chi \) the latter are generally weak, since in this compressed region only a small fraction of events pass the tight cuts applied on missing energy. Consequently, the dileptonic limits for spin-1/2 DM are seen to generally surpass the jets + MET limits. This is true of \( p\text{CS}_{RR}^u \) as well, except the bound on the coupling now rapidly tightens at \( m_\chi \approx 330 \text{ GeV} \). This happens because the production rate of the fermionic mediator (in \( p\text{CS}_{RR}^u, p\text{CS}_{RL}^u \)) is higher than the scalar mediator (in \( p\text{D}_{RR}^u, p\text{D}_{RL}^u \)), and we are able to saturate the CMS exclusion cross section with pure QCD production \( (\lambda \to 0) \) in this region. As for \( p\text{CS}_{RR}^d \), the dilepton bounds are weaker due to small down quark PDFs. At \( m_\phi = 2m_\chi \) the jets + MET limits are comparable to those at \( m_\phi = 1.1m_\chi \). This is because, though the signal MET acceptance improves in the uncompressed region, the mediator production rates fall with \( m_\phi \). As the dilepton signal is diminished in this region, it complements jets + MET in a model-dependent fashion: for spin-1/2 DM, the \( m_{ee} (\cos \theta_{CS}) \) bound outdoes jets + MET for \( p\text{D}_{RR}^u \) (\( p\text{D}_{RR}^d \)), while for spin-0 DM, no dilepton bound surpasses jets + MET.

The relic density constraint on DM overproduction is generally stronger for spin-0 DM than spin-1/2 DM. This is because the \( s \)-wave piece of complex scalar DM annihilation is chirality suppressed \[24, 40\]. At \( m_\phi = 1.1 \, m_\chi \) and \( m_\chi \leq 450 \text{ GeV} \), spin-0 DM gives weaker bounds since the efficient self-annihilation of the colored fermion mediator drives the co-annihilation mechanism in this region.

The dilepton probes greatly complement direct detection searches as well. Limits from the latter are generally strong when the mediator is near-degenerate with DM in mass, as seen in the \( m_\phi = 1.1 m_\chi \) plots. This is due to the factor of \( (m_\phi^2 - m_\chi^2)^{-k} \) in the
cross sections, where \( k = 4 \) (2) for spin-1/2 (spin-0) DM. The limit on spin-1/2 DM is mostly insensitive to \( \lambda \) at \( m_\chi \simeq 600 \text{ GeV} \) due to our scaling of the scattering cross sections with \( \Omega_\chi h^2/0.12 \propto \langle \sigma v \rangle \) – the former \( \propto \lambda^4 \) and the latter \( \propto \lambda^{-4} \) at large \( \lambda \), where co-annihilations with the mediators are unimportant. The limit does vary with the coupling at small \( \lambda \), where co-annihilations dominate. This asymptotic behavior of the limit with respect to \( \lambda \) allows our dilepton probes to constrain our set-up better than direct detection at \( m_\chi \gtrsim 600 \text{ GeV} \). On the other hand, the direct detection limit does not asymptote as quickly for spin-0 DM. This is because DM annihilations are chirality-suppressed in the \( s \)-wave, allowing co-annihilations to influence freeze-out even at large \( \lambda \). As a result, direct detection limits dwarf all other constraints for spin-0 DM at \( m_\phi = 1.1m_\chi \). The potency of dilepton probes is better at higher \( m_\phi \).

Due to the \( m_\phi^{-k} \) scaling, the limits weaken with \( m_\phi \) so much as to disappear from the \( m_\phi = 2m_\chi \) plots, allowing dilepons to probe this region better.

### 3.4.3 DM direct detection formulae

The spin-independent per-nucleon scattering cross sections, \( \sigma_{SI} \), is provided below for DD when DM behaves like a Majorana or real scalar particle:

- **Majorana DM** \((pD_{RR}, pD_{RL}, pD_{DR}, pD_{DL})\)

We have

\[
\sigma_{SI} = \frac{4}{\pi} \mu_{\chi N}^2 |f_N|^2, \tag{3.13}
\]

where \( \mu_{\chi N} \) is the DM-nucleon \((N = p, n)\) reduced mass, and effective coupling \( f_N \) is given by \([40]\)

\[
f_N \frac{m_N}{m_\phi} = f_q f_{T_u} + \frac{3}{4} (q_2 + \bar{q}_2) g_q - \frac{8\pi}{9\alpha_s} f_{T_G} f_G, \tag{3.14}
\]
with the Wilson coefficients

\[ f_q = \frac{m_\chi \lambda^2}{8(m_\phi^2 - m_\chi^2)^2}, \]
\[ f_G = \frac{-\alpha_s m_\chi \lambda^2}{96\pi m_\phi^2(m_\phi^2 - m_\chi^2)^2}, \] (3.15)

and the coefficients \( f_{T_u}(\text{proton}) = 0.023, f_{T_u}(\text{neutron}) = 0.017, f_{T_d}(\text{proton}) = 0.032, f_{T_d}(\text{neutron}) = 0.041, u_2 = 0.22, \bar{u}_2 = 0.034, d_2 = 0.11, \bar{d}_2 = 0.036, g_q = 4f_q, f_{T_G}(\text{proton}) = 0.925, f_{T_G}(\text{neutron}) = 0.922 \) \cite{75, 76}.

- Real Scalar DM (\( p_{\text{CS}}^{u_R}, p_{\text{CS}}^{u_L}, p_{\text{CS}}^{d_R}, p_{\text{CS}}^{d_L} \))

Here

\[ \sigma_{SI} = \frac{\mu_N^2}{\pi} \left( \frac{f_N}{m_\chi} \right)^2, \] (3.16)

with

\[ \frac{f_N}{m_N} = f_q f_{T_q} + \frac{3}{4}(q_2 + \bar{q}_2)g_q. \] (3.17)

The coefficients on the right-hand side are the same as before.
CHAPTER 4
NEW PHYSICS EFFECTS IN DIBOSON FINAL STATES

The tree-level partonic angular distribution of $W\gamma$ production possesses a feature known as the Radiation Amplitude Zero (RAZ). At the proton level the RAZ disappears, however, one can find a dip in the central region of the angular distributions, here called the Radiation Valley (RV). We show how searches for $W\gamma$ resonances get significantly improved if one focuses at the RV. We study how the significance improves in searches for scalar and vector resonances. Depending on the mass of the resonance, our results show an improvement of up to 65% in the significance for scalar resonances, and up to 100% for vector resonances. For the models studied, we also discuss the range of masses that can be probed at the high luminosity stage of the Large Hadron Collider (HL-LHC).

4.1 Motivation

Resonances are the most direct evidence of the existence of new degrees of freedom. The LHC searches for heavy resonances decaying to a photon and a hadronically decaying $Z/W/H$ boson. The results in [77] were obtained with the 13 TeV and 36.1fb$^{-1}$ data set. Minor excesses of 2.2 and 2.4 sigmas can be seen at masses of 1.7 and 2.6 TeV. Even though these are more than likely just statistical fluctuation, it is worth thinking about new physics models that can generate $W\gamma$ resonances. Some examples can be found in Quirk bounds states [78-80], compositeness [81-83], heavy

1Reference [111] provides a comprehensive summary of diboson searches including models and LHC searches.
vector triplets [84, 85], charged higgses [86, 87]. This chapter shows how to improve the sensitivity of $W\gamma$ resonances using the Radiation Amplitude Zero.

The RAZ is a spectacular feature of the partonic $W\gamma$ differential cross section $d\sigma/d\cos\theta^*$. It is an exact zero (at tree level) in the cross section at $\cos\theta^* = \pm 1/3$, where $\theta^*$ is the angle between the incoming parton and the outgoing photon, in the center of mass (CM) frame. The negative (positive) sign corresponds to $W^+$ ($W^-$) production. This phenomenon was discovered long ago by Brown, Mikaelian, Sahdev, and Samuel [88, 89] and it was explained by Brodsky et al. [90, 91] as the quantum version of the classical result that there is no dipole radiation in the scattering of particles with the same $e/m$ ratio. The classical relation $Q_1/m_1 = Q_2/m_2$ turns at the quantum level into $Q_1/(p_1 \cdot q) = Q_2/(p_2 \cdot q)$ [92], where $q_{1,2}$ are the charges of the colliding particles, $m_{1,2}$ their masses, $p_{1,2}$ their four-momentum, and $q$ is the four momentum of the outgoing photon. This formula defines the kinematic condition for which the amplitude of $W\gamma$ production is exactly zero at tree-level.

In going from parton level to a realistic hadronic collision, the zero gets washed out somewhat [93–101]. The largest contaminants are photon final state radiation (FSR), next-to-leading order corrections (NLO), and the reconstruction of the partonic center of mass frame. However, even including these washout effects, a clear dip in the angular distribution remains [102, 103]. In fact, evidence for the RAZ has been found in measurements of $\Delta y = y_\gamma - y_\ell$ at the Fermilab Tevatron in the $p\bar{p} \rightarrow \ell\nu\gamma$ channel [104]. More recently the CMS collaboration has confirmed these results measuring the charged rapidity difference [105, 106].

The goal of this chapter is to utilize the RAZ feature in $W\gamma$ to improve resonance searches, specifically, by focusing the searches at the dip of the angular distributions. Because the RAZ is an angular region where a small number of events are expected in the Standard Model (SM), using it as a kinematic cut improves the sensitivity of NP searches. Our results show how different angular distributions can be used to
Figure 4.1. Tree level contributions to $W\gamma$ production ($s$, $t$, and $u$ channels respectively). For leptonically decaying $W$ one should add an extra diagram where the photon is emitted by the lepton which corresponds to final state radiation.

improve the signal significance in the context of scalar and vector $W\gamma$ resonances. We then conclude by showing the mass range of resonances that can be probed at the LHC in the context of few models.

4.2 Radiation Amplitude Zero

At tree level, $W\gamma$ production occurs through the processes in Fig. 4.1. As one can see in Fig. 4.2 the parton level angular distribution shows an exact zero at $\cos \theta^* = \mp 1/3$ for $W^\pm$ production, where we have defined $\theta^*$ as the angle between the incoming quark and the photon in the CM frame. This zero occurs due to a cancellation between the diagrams in Fig. 4.1. Any NP signal can modify the shape of the angular distribution by spoiling such cancellation. If the NP effects come from higher dimensional operators or anomalous triple gauge couplings, the NP amplitude can interfere the SM amplitudes and reshape the angular distribution in nontrivial ways. If the NP effects come from either scalar or vector resonances, the interference effects are subdominant, but the contribution from NP can populate the dip in the angular distributions where the SM contribution is zero.

At hadron colliders, it is not possible to reconstruct the CM angle $\theta^*$ because we do not know the direction of the incoming quarks. However, by measuring the four-momentum of the final state photon and lepton it is possible to measure the rapidity
Figure 4.2. Partonic angular distribution for $W^\pm \gamma$ production. The RAZ is located at $\cos \theta^* = \mp 1/3$, where $\theta^*$ is the angle between the incoming quark and the photon in the CM frame.

distributions $y_\gamma$, $y_\ell$ and $\Delta y$. Also, it is possible to reconstruct the four-momentum of the final state neutrino by assuming that the leptons are produced by an on-shell $W$ (leptonic $W$ reconstruction). By knowing the momentum of all the external particles one can measure the CS angle. Here, the CS angle is calculated using Eq. 2.3, but this time $Q$ is the net momentum of the $W\gamma$ system with $Q_z$ ($Q_T$) the longitudinal (transverse) piece, and $p_i^\pm = (p_i^0 \pm p_i^z)/\sqrt{2}$, defining $p_1$ ($p_2$) as the momentum of the photon ($W$ boson). One of the goals of this analysis is to find evidences of the RAZ using these angular distributions; rapidities and the CS angle.

As mentioned before, one does not expect an exact zero in the hadronic distributions due to the smearing effects, mainly, the NLO corrections and the photon FSR contribution. Both effects imply to add extra diagrams to the ones in Fig. ref {diagrams}, which spoils the tree-level cancellation. The impact on the RAZ from NLO effects is minimized by imposing a jet veto on the events, whereas the impact due to the FSR contribution is minimized with cuts on the angle between the electron and the photon, $\Delta R_{e\gamma}$, and the photon isolation cone, $R_0$. Additional cuts on the invariant mass $m_{W\gamma}$ or the transverse cluster mass $m_T^{\ell \nu \gamma}$ of the events further
minimize the smearing effects. The variable $m_{T}^{\ell\nu\gamma}$ is defined as in [107] by

$$(m_{T}^{\ell\nu\gamma})^{2} = \left(\sqrt{m_{\ell\nu}^{2} + |p_{T}^{\ell} + p_{T}^{\nu}|^{2} + \text{MET}}\right)^{2} - |p_{T}^{\ell} + p_{T}^{\nu} + p_{T}^{\gamma}|^{2}, \quad (4.1)$$

where $m_{\ell\nu}$ is the invariant mass of the lepton-photon pair.

In Fig. 4.3 we show the SM expectation of different distributions that can be measured at the LHC in $W^{+}(\ell^{+}\nu)\gamma$ production. These distributions have been calculated at NLO including FSR effects using MCFM 8.0 [108–110], with the basic kinematic cuts

$$p_{T}^{\gamma} \geq 50 \text{ GeV}, \quad |y_{\gamma}| \leq 2.5,$$
$$p_{T}^{\ell} \geq 30 \text{ GeV}, \quad |y_{\ell}| \leq 2.5,$$
$$\text{MET} \geq 30 \text{ GeV}. \quad (4.2)$$

To suppress the NLO smearing we have used the jet veto

$$p_{T}^{j} \geq 30 \text{ GeV}, \quad |y_{j}| \leq 2.8, \quad (4.3)$$

and for the FSR smearing we have used the cuts

$$\Delta R_{\ell\gamma} \geq 0.7, \quad R_{0} \geq 0.4. \quad (4.4)$$

In Fig. 4.3, solid lines represent the distributions with only the basic kinematic cuts in Eqs. 4.2 4.3 4.4, while dashed lines include an extra cut in $m_{T}^{\ell\nu\gamma}$, or $m_{W\gamma}$, as indicated. These distributions in Fig. 4.3 show the following features:

- The difference in rapidities $\Delta y$ (Fig. 4.3, top) shows a dip in the central region related to the RAZ. For large values of the transverse mass, the depth of the central dip is bigger.
- The photon rapidity $y_{\gamma}$ (Fig. 4.3, center) is nearly flat for relatively low values
Figure 4.3. Proton level distributions of $W^+(ℓ^+ν)γ$ production: Difference of the photon and lepton rapidities (top), photon rapidity (center), and Collins-Soper angle (bottom). Solid lines represent the distributions with basic kinematic cuts (Eqs. 4.2, 4.3), while dashed lines include a cut in the cluster transverse mass $m_{ℓνγ}^T$ (Eq. 4.1), or the invariant mass $m_{Wγ}$, as indicated.

- The CS angle (Fig. 4.3, bottom) shows a minimum (or valley) at the center-left angular region $\cos θ_{CS} ∈ [−0.4, 0.1]$. This feature is a clear indicative of the underlying RAZ.

One can see how for large values of $m_{ℓνγ}^T$ or $m_{Wγ}$, the distributions in Fig. 4.3 are pushed towards the edges. Such behavior is characteristic of a $t$-channel process, in which the cross section does not fall as fast with energy at the edges due to the forward/backward divergences (also seen in Fig. 4.2 for $\cos θ^* = ±1$). Thanks to this
behavior, we see how the depth of the central region, here called Radiation Valley (RV), becomes deeper as the energy grows. This means that the RV is potentially a sweet kinematic spot to search for new physics; any NP signal whose angular distributions are expected to populate the central region can potentially be resolved by focusing the searches at the RV.

To quantify the depth of the RV, we define the ratio

$$R_V = \frac{\text{number of events inside the valley of } d\sigma}{\text{total number of events in } d\sigma},$$

where $d\sigma$ represents one of the distributions discussed above $\Delta y$, $y_\gamma$, and/or $\cos \theta_{CS}$. The valley is defined by a number $n$ of bins in the central region. Here we define $n$ such that the significance $S/\sqrt{B}$ is maximized\(^2\) for each distribution. The value of $n$ clearly depends on the particular search that the LHC is performing, for example, it is found that for scalar resonances $n = 8$ provides the highest increase in signal significance for the $\Delta y$ distribution. The variable $R_V$ is helpful to discriminate signals due to $W\gamma$ resonances from the SM background. The search strategy in the following:

- One first group the $\ell\gamma$+MET events into bins of $m_T^{\ell\nu\gamma}$ (or $m_{W\gamma}$ if leptonic $W$ reconstruction is applied).
- Using the angular distributions in each bin one then calculates the $R_V$ variable for both the background and the signal.
- Calculate the signal significance after applying the previous cuts.

The results show that the larger $m_T^{\ell\nu\gamma}$ (or $m_{W\gamma}$) the smaller $R_V$ is for the background, whereas $R_V$ increases for the signals, which means that the signal significance $S/\sqrt{B}$ also increases. This is illustrated in Fig. 4.4 where the data is split into invariant mass bins. The blue histogram represents the SM background and colored peaks represent different scalar resonances from the model introduced in Eq. 4.6 \(^2\) Where $S$ and $B$ represent the number of signal and background events respectively.
Figure 4.4. Invariant mass distribution of $W^+(\ell^+\nu)\gamma$ production. The blue histogram represents the SM background and colored peaks represent different benchmark signals from a scalar resonance. Solid lines represent full distributions whereas dashed lines are the distributions including a cut in the CS angle for each bin. For the 700 GeV bin, the number of events for both signal and background are explicitly shown by the little numbers next to the bins.

One can see for example, for the 700 TeV resonance, the significance increases from $280/\sqrt{140} = 24$ (solid lines), to $102/\sqrt{6} = 42$ (dashed lines). The distributions in Fig. 4.4 only include leading order calculations to illustrate the power of the technique here proposed. The results on the next sections will include NLO distributions.

The previous discussion is the main result of this paper and more details will be discussed below. We now want to stress out that while the RAZ is a tree-level result, the RV is the hadron-level remnant of it. From now on, the RV will refer to the dip in the central region of the angular distributions.
Figure 4.5. Production of $W\gamma$ though a charged squirk bound state. A s-up and a s-down quark are produced through a $W$ boson. They confine due to the dark confining force. This bound state is created with an arbitrary configuration of angular momentum $l$. The bound state relaxes very promptly and decays as a scalar resonance.

4.3 New Physics Signatures in $W\gamma$

Before studying the utility of the RV as a region in which to study new physics, it is important to discuss different models that can produce $W\gamma$ resonances. Here is provided an UV-motivated model, namely Squirks in Folded Supersymmetry [78, 79]. Also, some effective Lagrangians for vector and scalar $W\gamma$ resonances are considered. These new particles can be the pions and rho mesons of a multi-TeV confining force [81–85, 111–115], for example.

- Squirks in Folded Supersymmetry: In the Minimal Supersymmetric Standard Model (MSSM) each superpartner is charged under the same symmetries as its corresponding SM particle. This is important in order to cancel the quadratic divergences of the Higgs mass parameter. On the other hand, in Folded Supersymmetry the quark superpartners are not charged under the SM $SU(3)_c$, but rather under a mirror (or dark) color $SU(3)_{c'}$. In the UV theory there is a $Z_2$ symmetry that relates both $SU(3)_c \leftrightarrow SU(3)_{c'}$ which guarantees the cancelation of the quadratic divergences in the IR [116]. Due to this $Z_2$, the confining scales of both $SU(3)$ forces are close to each other. This implies that heavy superpartners tend to form bound states under the $SU(3)_{c'}$ force in the form of Quirks [117–122], or more accurately, Squirks (scalar quirks). If $\tilde{q}$ represents a quark superpartner and $\langle \tilde{q}\tilde{q}^* \rangle^Q$ is a squirk state with charge $Q$, one can imagine
these states to be produced at the LHC through Drell-Yan

\[ pp \to \tilde{q}\tilde{q}^* \to (\text{after confining}) \to \langle \tilde{q}\tilde{q}^* \rangle^Q. \]

If one produces a neutral squirk, it will mostly decay to dark glueballs. However, a charged squirk mostly decays to \( W\gamma \) as shown by the authors in [78]. The process is illustrated in Fig. 4.5, where one sees the parton level process through which a s-up and a s-down quarks are produced through a \( W \) boson. These squarks confine due to the dark force and form a charged squirk. This bound state is created with an arbitrary configuration of angular momentum \( l \). The system radiates away a large number of soft photons causing a very prompt relaxation. Once the system relaxes, it decays as a scalar resonance [78].

- **Effective heavy pions:** In composite models, it is possible to realize effective interactions involving heavy scalars with the SM gauge bosons. These anomaly-induced interactions have the form \( (1/\Lambda)\phi^a W^a_{\mu\nu} \tilde{B}^{\mu\nu} \), where \( \Lambda \) is related to the scale at which the UV theory confines. Here \( \phi^a \) is a pseudo scalar \( SU(2)_L \) triplet. We consider this interaction along with a coupling to matter \( y_m \) so that the resonance is produced through \( q'\bar{q} \) collisions

\[
\mathcal{L} \ni -\frac{g_m}{\Lambda} (iQ^t a H d^{\dagger} \phi^a + \text{h.c.}) + \frac{1}{\Lambda} \phi^a W^a_{\mu\nu} \tilde{B}^{\mu\nu}. \quad (4.6)
\]

- **Heavy Vector Triplets:** Besides heavy pions, a confining force can produce vector rho-like states. These heavy vectors can couple directly to matter through a renormalizable interaction, and to gauge bosons through higher dimensional operators. We consider the following interactions\(^3\) for our analysis

\[
\mathcal{L} \ni -g_m Q^t q^a \sigma^\mu Q V^\alpha_{\mu} + \frac{C_W}{\Lambda^2} V^a_{\mu\nu} W^a_{\nu\alpha} B_{\alpha} + \frac{C_h}{\Lambda^2} V^a_{\mu\nu} B^\mu B^\nu H^t a D^\dagger_{\nu} H. \quad (4.7)
\]

\(^3\)The interactions in 4.6 and 4.7 are CP invariant. The triplets \( \phi^a, V_m u^a \) can in principle mix with some of the SM particles. We assume that the effect of such mixing is small enough to be ignored.
where $g_m$ is the coupling to matter and the operator $c_W$ and $c_h$ provide two examples of mechanisms through which the heavy vector can decay to gauge bosons.

The previous interactions 4.6 and 4.7 produce $W\gamma$ resonances as illustrated in Fig. 4.6. For a scalar resonance both vertices are effective vertices whereas for vector resonances, the coupling to matter can be at tree-level.

4.4 Results

Fig. 4.7 shows the shape of signal distributions for $W^+(\ell^+\nu\gamma)$ resonances. Few values of the resonance mass $m_X$ has been used for illustration. The difference in
Figure 4.8. Ratio variable (left) defined in 4.5 that indicates how much the background decreases by making a cut in the central region of the rapidity distributions. For a given $R_V$, provided the signal shape (Fig. 4.7) this figure shows how much the signal significance (right) increases.

rapidities $\Delta y$ tends to populate the central bins. It presents a similar shape independently of the spin of the resonance. Another result, not shown in the figure, is that the shape of the $\Delta y$ signal does not vary much as the mass of the resonances change. The shape of the photon rapidity $y_\gamma$ tends to be slightly flat if the mass of the resonance is below 0.5 TeV. For large masses, the signal tends to accumulate in the central region (this result is independent of the spin of the resonance).

4.4.1 Rapidity distributions

One can highlight the following three results

(i) The signal histograms for the rapidity distributions are focused at the central regions,
(ii) The SM background histograms tend to populate the edges of the rapidity distributions, with few expected events at the RV,

(iii) The higher the energy, the deeper the RV.

These results indicate how powerful RV can be to discriminate the signal over the SM background. This point can be stressed out by the results in Fig. 4.8. The variable $R_V$ (dots) indicates how much the background gets reduced if one makes a cut in the central region ($R_V$ should be read from the scale to the left). Fig. 4.8 also shows how much the signal significance (lines) increases by focusing at the RV ($\Delta$Significance should be read from the scale to the right axis). In all cases (scalar and vector resonances), the difference in rapidities $\Delta y$ provides the highest increase in the signal significance at low masses, whereas $y_{\gamma}$ does it for large masses.

4.4.2 The role of the CS angle

As expected, the CS angle can be used to discriminate the spin of NP signals [20, 65]. This analysis becomes relevant once the LHC discovers a $W\gamma$ resonance; scalar resonances decay isotropically, which makes $\cos \theta_{CS}$ to be flat (except at the edges where the kinematic cuts of the final states objects reduces the number of events by a few). On the other hand, vector resonances tend to either populate the edges or the central region, depending on the type of interaction that one considers between the new particle and the vector bosons. For example, the operator $c_h$ tends to produce the longitudinal mode of the $W$ boson. This pushes the photon towards the edges of the distribution to compensate for the initial angular momentum provided by the pure left-handed interaction between the resonance and the quarks. The operator $c_W$ produces mostly the transverse components of the $W$ boson. This adds an extra vector to the final state that can compensate for the angular momentum of the initial states. This combination of spin and orbital angular momentum results in the angular pattern shown in Fig. 4.7, where the signal tends to populate the central region.
4.4.3 Mass sensitivity

The range of masses that can be probed with future luminosity is presented in Fig. 4.9. The figure shows the bounds on scalar resonances for the models discussed above. Similar results are expected for the vector resonances. The red line in Fig. 4.9 corresponds to the production of squirks bound states in Folded Supersymmetry. Few assumptions made are: 1. the bound state decays into $W\gamma$ 85% of the time following [78], and 2. only first and second generation of squirks contribute. (The third generation is expected to beta-decay before forming bound states.) This model is quite predictive in that the number of free parameters is small. The blue line represents a benchmark point in the parameter space of the Effective Heavy Pions model. We have taken a confining scale of $\Lambda = 10$ TeV, and a coupling to matter of $y_M = 0.14$. Solid lines represent the cross sections and expected sensitivity without using any cut in the rapidity distributions. The dashed lines is the analysis when one
focuses in the RV. The expected bounds on the masses of these resonances is in the range \(1.4 - 1.7\) TeV. More importantly, the bounds can be improved by \(200\) GeV by applying the cuts on the rapidity distributions here studied. This increase in mass corresponds to almost doubling the luminosity of the HL-LHC.
In the following, the summary and conclusions from chapters 3 and 4 are presented as well as a general conclusion on the thesis.

5.1 New Physics Signals in Dileptons

This chapter showed how dark matter may be characterized using invariant mass and scattering angle spectra of dilepton distributions at the LHC. If DM coupled to both quarks and leptons through $t$-channel mediators, radiative corrections from this dark sector, in combination with threshold effects, produce unique spectral features that may single out DM properties. The findings can be summarized:

1. Finding a dileptonic signal as sketched in Fig. 3.2 would imply that DM is not its own anti-particle.

2. The spin of DM is determinable from the $m_{\ell\ell}$ spectrum. In the region where the signal cross section rises quickly, its slope must be inspected; at higher $m_{\ell\ell}$, one must check whether the event ratio of signal over background grows rapidly or settles to a steady value.

3. If DM is a fermion, the signal cross section rises abruptly near twice the DM mass, thus revealing the mass of DM. This feature must be reflected as an abrupt change of slope at the same $m_{\ell\ell}$ in angular asymmetries, such as the forward-backward ($A_{\text{FB}}$) or center-edge asymmetry ($A_{\text{CE}}$).

4. The angular spectrum pinpoints the chirality of the fermions to which DM couples, an effect best seen in the angular asymmetries $A_{\text{FB}}$ and $A_{\text{CE}}$.

Having analyzed the signal features, constraints has been placed on the couplings and masses using LHC dilepton data, and contrasted them against bounds from multijets
+ MET searches, relic density measurements, and direct detection. One of the main results is that angular distributions provide better constraints than $m_{\ell\ell}$ distributions in some cases, significantly updating the conclusions of Ref. [39]. One can also see that dileptonic measurements in general are complementary to conventional DM searches. This is especially true of DM coupling to up quarks, where these probes often set the strongest collider bound on the RR model. It must be remembered that, just like jets + MET, the Drell-Yan process sets bounds not only on DM but on any analogous neutral particle that lives longer than collider time scales.

For the sake of illustration, only DM couplings to right-handed quarks have been considered. Should one of the dilepton signals here presented arise in forthcoming runs of the LHC, one must also entertain interpretations of DM coupling to other flavors and chiralities of quarks, and consider a wider range of splittings $\delta m$ when shape-fitting. Disentangling the exact Lagrangian structure would be a challenging task, and may involve deeper scrutiny of all available spectral information.

All in all, one can look forward to the amusing prospect of visible particle production educating us on dark matter.

5.2 New Physics Signals in Dibosons

Chapter 4 presented a discussion about the Radiation Amplitude Zero in $W\gamma$ production and the role it can play in detecting exotic resonances. For this, different kinematic distributions where the RAZ can be seen were presented, namely rapidity and Collins-Soper distributions. Some new physics models that produce $W\gamma$ resonances were provided as benchmark models. These effective interaction involve scalar and two cases of vector $W\gamma$ resonances; in one case the vector resonance produces mostly longitudinal $W$ bosons, and in the other case it produces mostly transverse $W$.

The main results of this chapter were the following: i) the SM background predicts
rapidity distributions that are spread towards the edges of the angular regions (proper
of a $t$-channel process), whereas the signal from the three type of resonances studied
tend to populate the bins in the central region of the angular distributions. For this
reason the rapidity distributions can be used as kinematic cuts to enhance the signal
significance of $W\gamma$ resonances. \textit{ii)} The shape of the Collins-Soper angular distribution
can be used to identify the spin of the resonance. If the resonance is a vector, $\cos \theta_{CS}$
can also identify if it produces mostly longitudinal or transverse $W$ bosons.

Each kinematic distribution here studied can play a role in the study of $W\gamma$ reso-
nances; the difference in rapidity $\Delta y$ provides the best sensitivity for small resonance
masses, the photon rapidity $y_\gamma$ does it for large masses, and the CS angle can be used
to discriminate the spin and couplings of the resonance.

5.3 A General Conclusion

As the LHC approaches to its era of high luminosity, it is important to identify
kinematic regions where the significance of new physics signals can be improved. This
is the spirit of the main two chapters of this thesis (Chaps. \[3\] and \[4\]). By measuring
angular distributions like the Collins-Soper angle and the rapidity distributions in
various final states, not only new resonances, but also different types of new physics
effects can potentially be resolved by their interference patterns. In such case, vari-
ables like $R_V$ and angular asymmetries like $A_{FB}$ and $A_{CE}$ can play an important
role. The thesis explored two specific examples where angular distributions can play
an important role unveiling new physics effects; in one case to discover Dark Matter
and to study its properties, and in the other example to explore exotic resonances in
$W\gamma$ production.
Note: During the time frame of this dissertation, other projects that ended up in publications were developed [123][129]. Here are summarized the general ideas and the main results in two of those projects.
APPENDIX A

THE HIGGS MASS IN SUPERSYMMETRIC EXTENSIONS OF THE SM

In order to reproduce the measured mass of the Higgs boson $m_h = 125$ GeV in the minimal supersymmetric standard model, one usually has to rely on heavy stops. By introducing a new gauge sector, the Higgs mass gets a tree-level contribution via a nondecoupling $D$-term, and $m_h = 125$ GeV can be obtained with lighter stops. This section studies the values of the stops masses needed to achieve the correct Higgs mass in a setup where the gauge group is extended by a single $U(1)_x$ interaction. The experimental limits on the mass of the $Z'$ gauge boson in this setup is presented, and then it is discussed how the stops masses vary as a function of the free parameters introduced by the new sector. The results of this section were published in [126].

A.1 Theoretical Set-up

The theoretical set-up is the same introduced by the authors of Ref. [130]. The MSSM superfields $\hat{L}, \hat{Q}, \hat{E}^c, \hat{D}^c, \hat{H}_u, \hat{U}^c$ and $\hat{H}_d$ are charged under the new gauge interaction as: 0, 0, 1/2, 1/2, -1/2 and -1/2. Right-handed neutrinos $\hat{N}$ have the opposite charge of the right-handed leptons. The model also introduces a pair of Higgs-like fields $\hat{\phi}$ and $\hat{\phi}^c$ with charges $\pm 1/2$ under $U(1)_x$. These fields are the responsible for breaking the new gauge symmetry. As discussed in [130], the radiatively corrected Higgs mass in the decoupling limit can be written as

$$m_h^2 = M_Z^2 \cos^2(2\beta) + \frac{1}{2} g_s^2 v^2 \cos^2(2\beta) \left(1 + \frac{M_{Z'}^2}{2 m_{\phi}^2}\right)^{-1} + \Delta m_{\text{loops}}^2,$$  \hspace{1cm} (A.1)
Figure A.1. Total cross section times branching ratio of the $Z'$ in the di-lepton channel as a function of the $Z'$ boson mass. The interception of the blue, orange and purple lines and the observed limit represent the lower bounds for the $Z'$ mass of our model for the given values of the gauge coupling $g_x$. The black line shows the limit on a Sequential Standard Model (SSM) $Z'$ [131], for comparison.

where the loop contributions depend (among other supersymmetric and soft parameters) on the stop mixing parameter $X_t = A_t - \mu \cot \beta$ and the supersymmetry breaking scale defined as

$$M^2_{SUSY} = m_{\tilde{t}_1} m_{\tilde{t}_2}. \quad (A.2)$$

A.2 Limits on the $Z'$ Mass

The $Z'$ in this model couples to SM fermions both through direct couplings and through $Z - Z'$ mixing. The expressions for the different couplings and mixing angle depend on the ratio $M_Z^2/M_{Z'}^2$, and can be found in [132] [133]. The decay of the $Z'$ is dominated by the channels $Z' \to f \bar{f}$, $Z' \to W^+W^-$ and $Z' \to Z h$. The partial widths in all modes grow linearly with $M_{Z'}$, so the branching ratios are basically constants in the limit when $M_{Z'} \gg M_Z$. 

68
This section follows a similar analysis to the one in [134–136] to set limits on new
gauge bosons using their decays to dileptons. For a given $Z'$ mass and $U(1)_x$ coupling,
one compares the cross section for production ($\sigma$) times branching ratio ($BR$) into
two leptons with experimental data. The left-hand panel of Fig. A.1 presents such
a comparison. The intersection of the observed limit (red line in Fig. A.1) and the
theory line for a particular value of $g_x$ gives the current lower $Z'$ mass bounds. For
$U(1)_x$ gauge couplings of $g_x = 0.32$, 0.63 and 0.84, there were found the limits of
$M_{Z'} \gtrsim 2.5$, 3.1 and 3.5 TeV, respectively using the run-I data. Because the $Z'$ can
decay to superparticles, one has to assume a spectrum in order to derive the limits.
Specifically, to generate Fig. A.1 the following naturalness-inspired spectrum was
considered: $\mu = 300$ GeV, $M_{\text{stops}} = 1$ TeV, $\tan\beta = 20$, and all other superpartner
masses at 4 TeV. With this setup, the only superpartners the $Z'$ can decay to are
third generation squarks and Higgsinos. It was found that the $Z'$ limits are fairly
insensitive to the masses of these states. Notice how the $g_x = 0.63$ exclusion is similar
to the limits on the Sequential Standard Model (SSM) [131]. This is expected as this
coupling is about the same size of the SM $g$ coupling and because the supersymmetric
channels do not contribute considerable to the width.

A.3 The Higgs Mass

As one can see in Eq. A.1, the Higgs mass contains two terms; the purely MSSM
piece and the $U(1)_x$ contribution. The MSSM piece is calculated using SUSYHD
[137], a public code that calculates the MSSM Higgs mass including two-loop thresh-
old corrections to the quartic Standard Model (SM) Higgs coupling. The code runs
down the SM parameters using three-loop renormalization group equations with a
leading four-loop QCD contribution to the strong gauge coupling. To calculate the
Higgs mass, the following parameters have been used: degenerate squarks ($m_{D_3}$ plus
first and second generations) and slepton masses $m_{\tilde{q}} = m_{\tilde{l}} = 4$ TeV, equal Wino
Figure A.2. Left (Right): Higgs mass as a function of $M_{\text{SUSY}}$ ($m_{\phi}$) for different values of the coupling $g_x$ and fixing $m_{\phi} = 1$ TeV ($M_{\text{SUSY}} = 700$ GeV). The solid lines correspond to the max-mixing case and the dashed lines to the non-mixing case. In both panels we have used the lower bound values for $M_{Z'}$ that correspond to the given $g_x$ values.

and Bino masses $M_2 = M_1 = 2$ TeV, gluino mass $m_{\tilde{g}} = 1.4$ TeV (close to the lower bounds reported by the LHC run I [138]) and a supersymmetric Higgsino mass of $\mu = 300$ GeV. Working with this spectrum, in order to reproduce the Higgs mass in the pure MSSM, one needs $M_{\text{SUSY}}$ to be around 1.6 TeV and 11.5 TeV in the cases of max-mixing ($\tilde{X}_t = X_t/M_{\text{SUSY}} = \sqrt{6}$) and non-mixing ($\tilde{X}_t = 0$), respectively. This behavior is consistent with previous work using different codes [139, 140]. Increasing the gaugino masses to 4 TeV causes a variation of less than 20% in $M_{\text{SUSY}}$.

Moving to the $U(1)_x$ extension, the picture changes, and the values of the Higgs mass are shown in Fig. A.2 for different values of the $U(1)_x$ coupling. The left panel shows the Higgs mass as a function of $M_{\text{SUSY}}$ with $m_{\phi}$ fixed to 1 TeV, while the right panel shows the Higgs mass for variable $m_{\phi}$ and fixed $M_{\text{SUSY}} = 750$ GeV (close to the highest lower bound reported for the stop mass [141–145]). Focusing on the left panel, for $m_{\phi} = 1$ TeV with $g_x = 0.32$, the new $D$-term gives a contribution to

\[^{1}\text{These extremum values of mixing are the ones that maximize and minimize the loop contribution from stops to the Higgs mass.}\]
Figure A.3. Masses of the stop eigenstates for large tan $\beta$, max-mixing with a small coupling $g_x = 0.32$ (solid lines) and non-mixing with a larger coupling $g_x = 0.63$ (dashed lines) cases. The gray-shaded region are excluded because $m_{\tilde{t}_2} > m_{\tilde{t}_1}$ and the gray-dotted contours correspond to fixed values of $M_{SUSY}$. The yellow region corresponds to a conservative bound that excludes the lightest stop for $m_{\tilde{t}_1} < 750$ GeV (refer to the text for details).

the Higgs mass such that the value of $M_{SUSY}$ reduces from 11.5 TeV to 8.1 TeV and from 1.6 TeV to 1.1 TeV for the no-mixing and max-mixing cases, respectively. For larger couplings one can see how the correct Higgs mass is achieved for even smaller values of $M_{SUSY}$, for example $g_x = 0.63$ implies $M_{SUSY} = 0.8$ TeV (max-mix) and $M_{SUSY} = 4.6$ TeV (non-mix), while $g_x = 0.84$ implies $M_{SUSY} = 0.6$ TeV (max-mix) and $M_{SUSY} = 3.3$ TeV (non-mix). Moving to the right panel, one sees that if $M_{SUSY}$ is fixed to 750 GeV, a $m_h = 125$ GeV cannot be generated in the no-mixing scenario with $g_x = 0.32$. For intermediate coupling, e.g. $g_x = 0.63$, the no-mixing scenario can be compatible with $m_h = 125$ GeV for $m_\phi$ around 3 TeV. For the max-mixing cases, small couplings $g_x = 0.32$ and $m_\phi$ around 2.6 TeV suffice.

A.4 Stops Masses

Using the definition (A.2) we can find the relation between the masses of the stop eigenstates corresponding to the contours showed in Fig. A.2. This draws a relation
between the new parameters $M_{Z'}$, $m_\phi$ and $g_x$ and the stops masses. This relationship is shown by the solid and dashed contours in Fig. A.3 for a fixed $m_\phi$. A given set of $U(1)_x$ parameters defines $M_{\text{SUSY}}$ and appears as a line in Fig. A.3. For example, for $g_x = 0.32$ and max-mixing, $m_\phi = 2.5$ TeV and the values $M_{Z'} = 3.1$, 3.5 and 4.0 TeV one generates the solid contours in the left panel of Fig. A.3 for which $M_{\text{SUSY}} = 830$, 880 and 940 GeV, respectively. Increasing the gauge coupling to $g_x = 0.63$, using the same $m_\phi$ and $M_{Z'}$ and assuming no stop mixing yields the dashed contours; the corresponding $M_{\text{SUSY}}$ values are 980, 1200 and 1500 GeV, respectively. A similar analysis is shown in the right panel of Fig. A.3 but considering a wider range of $M_{Z'}$ and $m_\phi$. Both panels show the same trend; increasing the gauge coupling or the soft term $m_\phi$ increases the D-term contribution, thereby causing the value of $M_{\text{SUSY}}$ to decrease and shifting the stop mass contours to lower values. This behavior can be seen by comparing the solid orange line in the left panel with the green solid line in the right panel; these lines correspond to $g_x = 0.32$, max-mixing, $M_{Z'} = 3.5$ TeV and $m_\phi = 2.5$ and 3.0 TeV, respectively. On the other hand, as one increases the mass of the $Z'$, the effects of the D-terms reduce, and $M_{\text{SUSY}}$ and the stop mass contours move to higher values. For example, increasing $M_{Z'}$ from 2.5 to 3.5 TeV changes $M_{\text{SUSY}}$ from 730 to 810 GeV, for $g_x = 0.32$ and max-mixing, and increasing $M_{Z'}$ from 3.1 to 4.0 TeV, $M_{\text{SUSY}}$ changes from 760 to 1100 GeV for $g_x = 0.63$ with non-mixing.

On the experimental bounds of the stops - Although there are a plethora of stop searches at the LHC, each involves some degree of assumption about the stop branching ratio, the mass difference between the stop and its decay products, and the composition of the stops and lightest superparticle (LSP) [141]-[145]. As a result, it is almost impossible to derive a model-independent bound. In the supersymmetric spectrum here used so far the LSP mass is 300 GeV, value for which there was no
bound on the lightest stop\footnote{Obviously, the mass of the lightest stop has to be always greater than $\mu$.} for the run-I of the LHC. Decreasing $\mu$ (and therefore the LSP mass) will lead to an limit between roughly 400 GeV to 750 GeV for $\tilde{t}_1$. In Fig. A.3 we show the 750 GeV limit as the lower bound for the lightest stop, although this should be viewed as a conservative bound -- most stops searches suppose a 100% BR of the stop into top plus neutralino, while in this case there is a sensible probability of stop decay into bottom quark plus chargino. In either decay mode, the bounds become looser if the spectrum is highly compressed since the stop decay products become soft and cannot be efficiently identified. As these bounds are independent on the details of the $U(1)_x$ sector, they will not change if one increases or decreases the mass of the $Z'$. 

A.5 Conclusions

This section studied the interplay of an extra $U(1)_x$ non-decoupling D-term contribution to the Higgs quartic and the stop masses needed to reproduce $m_h = 125$ GeV. It was found that the correct Higgs mass can be reproduced with stops in a region between 700 – 800 GeV and a $Z'$ resonance close to the 2.5 TeV bound from the run I of the LHC, or in a higher region 800 – 900 GeV if the $Z'$ resonance is heavier (3.1 TeV). This represents a scenario more optimistic than the pure MSSM where the stops masses are expected to be at the multi-TeV range, out of the reach of the LHC.

A general conclusion for models where the Higgs mass is a prediction, like in supersymmetry, is that the measured value of 125 GeV can put both lower and upper bounds on the masses of new particles.
APPENDIX B

LEPTOQUARKS AND THE R(D) ANOMALY

Over the last few years there has been a persistent disagreement between the SM prediction and experimental measurements of

\[ R_D = \frac{B(\bar{B} \to D\tau\bar{\nu})}{B(\bar{B} \to D\ell\bar{\nu})}, \quad R_{D^*} = \frac{B(\bar{B} \to D^{*}\tau\bar{\nu})}{B(\bar{B} \to D^{*}\ell\bar{\nu})}, \]

where \( \ell = e, \mu \). The SM predictions \([146-153]\) are known at the 1−2\% level

\[ R_D(\text{pred}) = 0.299 \pm 0.003, \quad R_{D^*}(\text{pred}) = 0.258 \pm 0.005, \]

whereas the combination of all current data shows \([154-161]\)

\[ R_D(\text{exp}) = 0.407 \pm 0.046, \quad R_{D^*}(\text{exp}) = 0.306 \pm 0.015, \]

which indicates a 2.5\( \sigma \) deviation of the data above the SM predictions. This anomaly may be addressed by introducing interactions beyond the Standard Model involving new states. Since the processes involved are quark flavor changing, any new states would need to couple to at least two different generations of quarks, requiring a non-trivial flavor structure in the quark sector while avoiding stringent constraints from FCNC processes. In this section, scalar leptoquarks (LQ) are introduced as a possible solution for the \( R_{D(\ast)} \) anomaly under the assumption of minimal flavor violation (MFV). Here have been considered all possible representations for the LQ under the SM quark flavor symmetry group, consistent with asymptotic freedom. Constraints
on the parameter space were derived from: self-consistency of the MFV scenario, perturbativity, the FCNC decay $b \to s\nu\bar{\nu}$ and precision electroweak observables. The main result is that none of the scalar LQ in the context of MFV can explain the $R_{D(\ast)}$ anomaly while simultaneously avoiding all constraints within this scenario. The results of this section were published in [123].

B.1 Theoretical Framework

LQ are hypothetical new particles that can interact with both leptons and quarks. A typical LQ interaction with SM fields looks like

$$L \ni \lambda_{ij} \phi \bar{q}_i \ell_j,$$  \hspace{1cm} (B.1)

where $\lambda_{ij}$ is the coupling, $\phi$ the LQ field, $q_i$, $\ell_j$ are any SM quark and lepton respectively, and $ij$ represent flavor indices.

One of the major naturalness problems that one faces with the introduction of LQ is their contribution to FCNC. Because the couplings $\lambda_{ij}$ are arbitrary, LQs can mediate new flavor-changing processes at rates far beyond those allowed by the SM. In order to have LQ anywhere near the weak scale, one must impose on the couplings some structure that, once the quark and lepton fields have been rotated to their mass eigenbasis, do not mix quarks of different generations at the tree level in neutral currents. Such a requirement would be highly unnatural unless the LQ interactions knew about the structure of Yukawa interactions and were somehow aligned with them. This is precisely what one finds in the Minimal Flavor Violation scheme.

In MFV, one promotes the approximate $SU(3)$ flavor symmetries of the quarks (and in some cases also the leptons) to exact symmetries that are broken only by the Yukawa interactions, which one treats as spurions. In such a case, the diagonalization of the quark masses similarly diagonalizes all other flavor symmetry-breaking
terms, up to corrections due to the off-diagonal elements of the CKM matrix. Before specifying the family of LQ models and the MFV interactions of interest, it is helpful to examine solutions to the $R_{D^(*)}$ anomaly within an effective Lagrangian approach. This will limit right away the number of LQ models that can potentially explain the anomaly.

B.2 Operator Analysis

Here, the fits found in Ref. [162] were updated with the most recent world averages from the Heavy Flavor Averaging Group (HFLAV). The pieces of the effective Hamiltonian contributing to $b \to c\tau\bar{\nu}_\tau$ can be written as:

$$\mathcal{H}_{eff} \ni 2\sqrt{2}G_F V_{cb}\mathcal{O}_{VL} + \frac{1}{\Lambda^2} \sum_i C_i \mathcal{O}_i,$$

where $\mathcal{O}_{VL}$ is the SM operator, $(\bar{c}\gamma_\mu P_L b)(\bar{\tau}\gamma^\mu P_L \nu_\tau)$, and $\mathcal{O}_i (C_i)$ are the dimension-6 Wilson operators (and their coefficients). The complete list of dimension-6 operators that can contribute to $\bar{B} \to D^{(*)}\tau\bar{\nu}_\tau$ is given in Ref. [162]. Of these operators, three can be mediated by scalar LQ. These three operators, denoted $\mathcal{O}'_{SL}$, $\mathcal{O}'^{''}_{SL}$ and $\mathcal{O}'^{''}_{SR}$, are listed in Table B.1 along with the SM quantum numbers of the corresponding scalar LQ. Of these three, the operator $\mathcal{O}'^{''}_{SR}$ is identical to the $V - A$ operator of the SM (up to a factor of 2) after a Fierz transformation of the fermion fields, the operator $\mathcal{O}'^{''}_{SL}$ becomes a combination of scalar and tensor operators.

A $\chi^2$ analysis can be done for each operator separately, using as inputs the current combined best fit values for $R_D$ and $R_{D^*}$ given in section above. The results of such a fit are shown in Fig. B.1 assuming that all three operators share roughly the

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1 https://hflav.web.cern.ch/

2 Throughout this analysis, it is assumed that the neutrino in the final state of $\bar{B} \to D^{(*)}\tau\bar{\nu}_\tau$ is the $\nu_\tau$ of the SM.
same systematics and efficiencies in their contributions to $R_{D^{(*)}}$. Here the scale $\Lambda$ is set to 1 TeV, and the value of $\chi^2$ is plotted as a function of the corresponding Wilson coefficient. One sees immediately that the operator $O_{S_L}''$ or $O_{S_R}''$ can provide a significantly improved fit to the $R_{D^{(*)}}$ data as compared to the SM (which is denoted with the solid line at $\chi^2 \simeq 15$). On the other hand, the operator $O_{S_L}'$ provides a fit to the data that is only minimally better than the SM in terms of total $\chi^2$, at the cost of an additional degree of freedom. One can therefore judge this operator as providing a poor explanation for the anomaly and it is not considered further.

Thus, there are only two scalar LQ that one needs to consider further: the $S_1$, an $SU(2)_L$ singlet with hypercharge of $1/3$; and $S_3$, an $SU(2)_L$ triplet with hypercharge also of $1/3$.

The value of the Wilson coefficients for which $\chi^2$ is minimized are

$$C_{S_L}'' = -0.428 \pm 0.096,$$  \hspace{1cm} (B.2)

and

$$C_{S_R}'' = 0.293 \pm 0.074 \quad \text{or} \quad -5.72 \pm 0.07,$$  \hspace{1cm} (B.3)
Figure B.1. $\chi^2$ values from a fit to the $R_{D^{(*)}}$ data, as a function of Wilson coefficients ($C_i$) for the operators generated by LQ exchange.

for $\Lambda = 1$ TeV. These values are indicative of the lengths to which one must go in order to solve the $R_{D^{(*)}}$ anomaly by the exchange of new particles. If one identifies the coefficients $C''$ with $\lambda^2/M_{LQ}^2$, with $\lambda$ representing the coupling of the LQ to the fermions, and $M_{LQ}$ its mass, then for $\lambda \sim O(g)$ one finds $M_{LQ} \sim 300$ GeV to 1.3 TeV.

B.3 Leptoquark Interactions and Flavor Structure

The most general Lagrangians for $S_1$ and $S_3$ that preserve both Baryon ($B$) and Lepton ($L$) numbers can be written as:

$$\mathcal{L}_{S_1} = S_1 \left\{ \lambda_{ij} \bar{Q}_i c_i i\tau_2 L_j + \bar{\lambda}_{ij} \bar{u}_i c_j e_j + h.c. \right\}, \quad (B.4)$$

and

$$\mathcal{L}_{S_3} = S_3^a \lambda_{ij} \bar{Q}_i c_i i\tau_a L_j + h.c. \quad (B.5)$$

In the expressions above, $i, j$ are generation indices, $Q_i$ and $L_j$ are the left-chiral quark and lepton doublets, $u_i$ and $e_j$ are the right-chiral quark and lepton singlets, and $\tau_a$
\[(a = 1, 2, 3)\) are the \(SU(2)\) generators. In order to impose MFV on the Lagrangians above, one must make several assumptions:

- Leptoquark interaction terms are required to be singlets under the SM flavor symmetry group in the quarks sector, i.e., \(G_q \equiv SU(3)_Q \times SU(3)_u \times SU(3)_d\). Doing this imposes a symmetry structure on the leptoquarks and coupling constants. The leptoquarks themselves must transform under \(G_q\) and the coupling constants must be written as an expansion in the Yukawa matrices \(Y_U\) and \(Y_D\), which transform under \(G_q\) as \((3, \bar{3}, 1)\) and \((3, 1, 3)\), respectively. Thus, these models contain a multiplicity of leptoquarks, only a few of which will be relevant for explaining the \(R_{D(\tau)}\) anomaly.

- Leptoquarks only couple to the \(\tau\) lepton, a requirement sometimes called “\(\tau\) alignment” in the literature. In principle, one could try to perform an MFV-like analysis in the lepton sector but it actually does not affect the results of this analysis. The present assumption indicates that the solution to the \(R_{D(\tau)}\) anomaly is entirely due to new physics contributions to the \(b \to c\tau\bar{\nu}_\tau\) process, with no new contributions to \(b \to c\ell\bar{\nu}_\ell\) for \(\ell = e, \mu\).

Under these assumptions, the Lagrangians above reduce to the following forms for \(S_1\) and \(S_3\),

\[
\mathcal{L}_{S_1} = \lambda_{1L}(S_1Y)^i\bar{Q}_i i\tau_2 L_3 + \lambda_{1R}(S_1Y')^k \bar{u}_k \tau + h.c.
\]

\[
\mathcal{L}_{S_3} = \lambda_3(S_3a Y'')^i\bar{Q}_i i\tau_2 \tau a L_3 + h.c.
\]

(B.6)

Here the \(\lambda_i (i = 1L, 1R, 3)\) are the overall couplings that multiply \(G_q\)-invariant terms; refered from now on as the “universal” couplings. \(L_3\) is the 3rd-generation lepton doublet, and \(\tau\) the right-handed charged tau lepton. The \(Q\) and \(u\) fields are taken to transform as fundamentals under their respective flavor groups \(SU(3)_Q\) and \(SU(3)_u\). The terms set off in parenthesis, such as \((S_1Y)^i\), are contractions of the leptoquark flavor multiplets with some number of Yukawa matrices (here collectively denoted \(Y\)) that serve as spurions under \(G_q\). The products \((S_1Y)\) and \((S_3Y'')\) transform as a \((\bar{3}, 1, 1)\) under \(G_q\), while \((S_1Y')\) transforms as a \((1, 3, 1)\). The contractions between the leptoquark field and the Yukawa spurions can be quite complicated; this analysis is not confined to the often-used simplification that \(Y\) transforms simply as an adjoint
of one of the three $SU(3)$ flavor groups. This allows one to probe flavor structures that are often ignored in MFV analyses.

In general, the $Y$ (or $Y', Y''$) term that appears in the equations above is a linear combination of an arbitrary number of product of Yukawa matrices:

$$Y \sim a_0 + a_1 Y_u + a_2 Y_u Y_u^\dagger + ... + b_2 Y_d Y_d^\dagger + ... + c_1 Y_u Y_d + ...$$

where some of the coefficients $a_i, b_i, c_i$ are zero depending on how the leptoquark transforms under $G_q$, and depending on whether contractions between the spurions, the leptoquark, and the quark doublet can generate $G_q$-invariant terms as required in the analysis.

Having now a form for the Lagrangian, one needs to identify all possible representations of $G_q$ under which the leptoquark fields can transform. However, in order to simplify the analysis, the following additional constraints are imposed:

- Because of the potentially large numbers of leptoquarks being introduced, all of which transform as triplets under QCD, asymptotic freedom could be lost in the theory. Therefore, it is reasonable to require the total number of QCD triplets not to exceed 16, at which point the one-loop QCD $\beta$-function flips sign. This requirement will limit the number of possible representations here considered, though it is not an absolute requirement for a self-consistent, low-energy theory.

Provided this, one finds that $S_1$ and $S_3$ can only have the following quantum numbers under $G_q$:

$$S_1 : (\bar{3}, 1, 1), (1, \bar{3}, 1), (1, 1, 3)$$

$$\phantom{S_1} : (6, 1, 1), (1, 6, 1), (1, 1, 6)$$

$$\phantom{S_1} : (3, 3, 1), (3, 1, 3), (1, 3, 3)$$

$$S_3 : (\bar{3}, 1, 1), (1, \bar{3}, 1), (1, 1, 3).$$
From this point on, these quantum numbers under $G_q$ will be referred to as flavor charges. In the analysis that follows, the Wilson coefficients $C_{S_L}''$ and $C_{S_R}''$ are calculated by integrating out LQ with any of the above flavor charges. In addition to $C_{S_L}''$ and $C_{S_R}''$, the Wilson coefficient for the operator $(s^c P_R \bar{\nu})(\bar{b}^c P_L \nu) \equiv O_{bs \nu \bar{\nu}}$ is calculated as it generates the decay $b \rightarrow s \nu \bar{\nu}$. Such operator is used later to set constraints on the MFV models here studied.

Finally, to guarantee the self-consistency of the MFV framework \[163\], it is required that

- The coefficients $a_i$, $b_i$, $c_i$ in the expansion of Eq. B.7 are $O(1)$ or smaller, whereas the overall couplings $\lambda_i$ in Eq. B.6 cannot be larger than $\sqrt{4\pi}$ to ensure that the interactions are perturbative.

### B.4 Constraints and Results

Eq. B.8 shows the 9 (3) choices of flavor charges for $S_1$ ($S_3$) that can explain the $R_{D^{(*)}}$ anomaly under the assumption of MFV. Constraints from the perturbativity of the universal couplings and from the meson decay $b \rightarrow s \nu \bar{\nu}$ rule out most of these models (the details are explained in \[123\]). Only two LQ are not disallowed by these two constraints: $S_1(3, 1, 1)$ and $S_1(1, 1, \bar{3})$. The next section focuses on the constraints on these two models due to EW observables.

The dominant effect of the $S_1$ leptoquark on precision electroweak observables is through its modification of the couplings of $Z$ to fermions at one loop; the relevant diagrams are shown in Fig. B.2. A number of electroweak observables can be impacted by the presence of a (predominantly) third-generation leptoquark, including the invisible width of the $Z$, the forward-backward asymmetry of $Z \rightarrow \bar{b}b$ or $\bar{\tau}\tau$, or the total rates for these same two processes. The results in this section show that the partial decay width of $Z \rightarrow \tau\bar{\tau}$ imposes the strongest constraints on $S_1(3, 1, 1)$ and $S_1(1, 1, \bar{3})$. These two flavor charges are considered below.
S\(_1(\overline{3}, 1, 1)\): The relevant part of the Lagrangian for the triplet of \(S_1(\overline{3}, 1, 1)\) leptoquarks is:

\[
\mathcal{L} \supset \lambda_{1L} S_1^{q_i} \left\{ (V^{+ \bar{u}_L^c})_{q_i} \tau_L - \bar{d}_L^{\tau_1} \nu_L \right\} \\
= \lambda_{1L} S_1^1 \left\{ (V_{ud} \bar{u}_L^c + V_{cd} \bar{c}_L^c + V_{td} \bar{t}_L^c) \tau_L - \bar{d}_L^{\tau_1} \nu_L \right\} \\
+ \lambda_{1L} S_1^2 \left\{ (V_{us} \bar{u}_L^c + V_{cs} \bar{c}_L^c + V_{ts} \bar{t}_L^c) \tau_L - \bar{s}_L^{\tau_1} \nu_L \right\} \\
+ \lambda_{1L} S_1^3 \left\{ (V_{ub} \bar{u}_L^c + V_{cb} \bar{c}_L^c + V_{tb} \bar{t}_L^c) \tau_L - \bar{b}_L^{\tau_1} \nu_L \right\}
\]

Note that among the flavor triplet of leptoquarks, only \(S_1^3\) contributes to \(R_{D^{\phi}}\), and can explain the anomaly for \(\lambda_{1L} \simeq 2.8\) (for \(M_{S_1} = 1\) TeV). Such large value of \(\lambda_{1L}\) is needed to make up for the \(V_{cb}\) suppression in the \(S_1^3\bar{c}_L^c \tau_L\) coupling. But this implies that the \(S_1^3\bar{t}_L^c \tau_L\) coupling, which is only “suppressed” by \(V_{tb}\), becomes fairly large. As it is shown below, this fact makes this particular LQ interaction (with a \(t\)-quark) very sensitive to electroweak observables.

Before calculating the constraints from electroweak precision measurements, note that the \(\tau\) lepton has a large coupling to the \(u\)-quark through two of the components
of the leptoquark flavor triplet: $S_1^1$ and $S_1^2$. These components can therefore be strongly constrained by $\bar{\tau}\tau$ production at the LHC. However, one can avoid such constraints, if needed, by assuming $S_1^1$ and $S_1^2$ to be much heavier than $S_1^3$, which is allowed within the framework of MFV, since the mass of $S_1^3$ can split from its flavor partners due to the large top Yukawa [164]. For completeness, we will calculate below the electroweak constraints both with and without decoupling the leptoquark components $S_1^1$ and $S_1^2$.

Assuming that the flavor triplet of leptoquarks is degenerate, then the shift in the coupling of the $Z$ to leptons can be expressed as:

$$
\Delta g^S_1(Z \to \bar{\tau}\tau) = \frac{3g_2|\lambda_{1L}|^2m_t^2}{32\pi^2 c_w M^2} \left(2 \log \frac{M}{m_t} - 1\right) + \frac{2g_2|\lambda_{1L}|^2m_Z^2}{96\pi^2 c_w M^2} \times \left\{ \left(-\frac{1}{2} + s_w^2\right) - \left(\frac{1}{2} - \frac{2}{3}s_w^2\right) \left(12 \log \frac{M}{m_Z} + 1 + i6\pi\right) \right\},
$$

(B.9)

$$
\Delta g^S_1(Z \to \bar{\nu}\nu) = \frac{3g_2|\lambda_{1L}|^2m_Z^2}{96\pi^2 c_w M^2} \times \left\{ \frac{1}{2} - \left(-\frac{1}{2} + \frac{1}{3}s_w^2\right) \left(12 \log \frac{M}{m_Z} + 1 + i6\pi\right) \right\}.
$$

(B.10)

For the $S_1(3, 1, 1)$ leptoquark, there is no equivalent correction to the right-handed couplings of the $Z$ to fermions due to the structure of the Lagrangian [B.6]. These relations are valid up to leading order in $m_Z/M$ and $m_t/M$ (where $M$ is the mass of $S_1$). In the expression above, the first term dominates and is due to the contribution of the top quark in the loop, where it picks up an enhancement $\propto m_t^2$ from helicity flips on each of the $t$-quark lines. The second, smaller contribution is due to $u$- and $c$-quarks running in the loop, and includes an imaginary component when the quarks go on shell. This smaller contribution has the opposite sign to the dominant term and is thus included in the calculations in order to obtain conservative bounds. At the
same time, the shift in the $Z$ coupling to $\bar{\nu}\nu$ is due to down-type quarks in the loop, and picks up no large enhancements. (Shifts in the $Z$ couplings to quarks involve only leptons in the loops and are even smaller.)

The LQ in this set-up predominantly modify the $Z$ decay to $\tau$, $\nu_\tau$, and $b$. The affected LEP and SLC $Z$-pole observables with their SM predictions and measured values are summarized in Table B.2 where $\Gamma_Z$ is the total decay width of $Z$, $\Gamma(\text{inv})$ is its invisible decay width, $R_\tau \equiv \Gamma(\text{had})/\Gamma(\tau\bar{\tau})$ and $R_b \equiv \Gamma(b\bar{b})/\Gamma(\text{had})$. $\Gamma(\text{had})$ is the partial width of $Z$ into hadrons, which receives a new contribution from the leptoquark-mediated $Z \to b\bar{b}$ process. The terms $A_{F_B}^{(0,f)}$ and $A_f$ ($f = \tau, b$) are the asymmetry observables that quantify parity violation in weak neutral currents. These are defined as:

$$A_f \equiv \frac{2g_A^f g_V^f}{(g_A^f)^2 + (g_V^f)^2}, \quad A_{F_B}^{(0,f)} = \frac{3}{4} A_e A_f,$$

where, $g_V^f$ and $g_A^f$ are the effective vector and axial couplings of $Z$ to fermions ($f = \tau, b$).

Figure B.3 summarizes the electroweak constraints on $S_1(\bar{3}, 1, 1)$. The green region is the parameter space where the $R_{D^0}$ anomaly can be explained by this choice of flavor charge, within either $1\sigma$ (or $2\sigma$) of the experimental measurements. Meanwhile, the constraints from the electroweak data disfavor the region above the dashed blue line in Fig. B.3 at 95% C.L. (i.e., $\Delta\chi^2 \geq 5.99$ for two free parameters). These constraints are obtained by a $\chi^2$ fit to all the electroweak observables that are strongly affected by the presence of leptoquarks. These observables include $R_\tau$, $\Gamma(\text{inv})$, $R_b$, $A_\tau$ and $A_b$ displayed in Table B.2. The current direct pair production bounds from LHC are also indicated by the red vertical line.

If one assumes that the leptoquarks $S_1^1$ and $S_1^2$ are much heavier than $S_1^3$, then the electroweak constraints become even stronger, with the new 95% C.L. now indicated by a solid blue line in Fig. B.3. This is because the contributions from $u$- and $c$-
TABLE B.2

RELEVANT LEP AND SLC OBSERVABLES

<table>
<thead>
<tr>
<th>Observable</th>
<th>Experimental</th>
<th>Standard Model</th>
<th>Pull</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\Gamma_Z$ [GeV]</td>
<td>$2.4952 \pm 0.0023$</td>
<td>$2.4943 \pm 0.0008$</td>
<td>0.4</td>
</tr>
<tr>
<td>$\Gamma_{(had)}$ [GeV]</td>
<td>$1.7444 \pm 0.0020$</td>
<td>$1.7420 \pm 0.0008$</td>
<td>-</td>
</tr>
<tr>
<td>$\Gamma_{(inv)}$ [MeV]</td>
<td>$499.0 \pm 1.5$</td>
<td>$501.66 \pm 0.05$</td>
<td>-</td>
</tr>
<tr>
<td>$R_\tau$</td>
<td>$20.764 \pm 0.045$</td>
<td>$20.779 \pm 0.010$</td>
<td>-0.3</td>
</tr>
<tr>
<td>$R_b$</td>
<td>$0.21629 \pm 0.00066$</td>
<td>$0.21579 \pm 0.00003$</td>
<td>0.8</td>
</tr>
<tr>
<td>$A_{FB}^{(0,\tau)}$</td>
<td>$0.0188 \pm 0.0017$</td>
<td>$0.01622 \pm 0.00009$</td>
<td>1.5</td>
</tr>
<tr>
<td>$A_{FB}^{(0,b)}$</td>
<td>$0.0992 \pm 0.0016$</td>
<td>$0.1031 \pm 0.0003$</td>
<td>-2.4</td>
</tr>
<tr>
<td>$A_\tau$</td>
<td>$0.1439 \pm 0.0043$</td>
<td>$0.1470 \pm 0.0004$</td>
<td>-0.7</td>
</tr>
<tr>
<td>$A_b$</td>
<td>$0.923 \pm 0.020$</td>
<td>$0.9347$</td>
<td>-0.6</td>
</tr>
</tbody>
</table>

quarks to $\Delta g_L^{S_1}(Z \rightarrow \bar{\tau}\tau)$ have the opposite sign to that from the $t$-quark and, after decoupling $S_1^1$ and $S_1^2$, those negative contributions become suppressed.

As Fig. 3 demonstrates, there is no parameter space remaining in which $S_1(\bar{3}, 1, 1)$ can explain the $R_{D^0}$ anomaly while evading electroweak constraints. This is true either in the case with decoupled $S_1^1$ and $S_1^2$, or without.

$S_1(1, 1, \bar{3})$: The contributions from these leptoquarks to $\mathcal{O}_{SR}^{S_1}$ mimic exactly those of the previous case, but now with $\lambda_{1L}$ replaced by $\lambda_{1L} y_b$. As such, they are highly suppressed for SM-like bottom Yukawa couplings, but could be sizable in a two-Higgs doublet model with large $\tan\beta$, where the bottom Yukawa can be $O(1)$. In either case, the leading term in the Lagrangian of $S_1(1, 1, \bar{3})$, can be written as:

$$\mathcal{L} \supset \lambda_{1L} y_b S_1^1 \{ (V_{ub}\bar{u}_L + V_{ub}\bar{c}_L + V_{tb}\bar{t}_L) \tau_L - \bar{b}_L \nu_L \} ,$$
Figure B.3. Parameter space of the leptoquark $S_1(3, 1, 1)$ in which it can explain the $R_{D^0}$ anomaly (green shaded region), along with the regions excluded by electroweak precision measurements (above dashed blue line) and pair production bounds (left of solid red line). The solid blue line indicates the electroweak constraints after decoupling the $S_1^{1,2}$.

ignoring terms suppressed by $y_d$ and $y_s$. For $y_b = 1$, this Lagrangian is exactly that of the $S_1^3$ component of $S_1(3, 1, 1)$. This implies that Fig. B.3 can also be used to study the parameter space of $S_1(1, 1, 3)$, with $|\lambda_{1L}|$ on the y-axis replaced by $|\lambda_{1L}y_b|$ and constraints from electroweak precision measurements denoted by the solid blue line. As before, there is no parameter space for $S_1(1, 1, 3)$ where it can both avoid the electroweak constraints and explain the $R_{D^0}$ anomaly.

B.5 Summary and Conclusions

This section studied the feasibility of scalar LQ as an explanation of the $R_{D^0}$ anomaly under the assumption of Minimal Flavor Violation (MFV). The family of scalar LQ models and the MFV set ups that can be written down is countless. To narrow down the models that can potentially explain the $R_{D^0}$ anomaly, the following
line of logic was followed:

- First, studying the effective operators that improve the theoretical prediction with respect to the data, it was found that only two scalar LQ, $S_1$ and $S_3$ are to be considered. At this point, all possible quantum numbers under the flavor group $G_q = SU(3)_Q \times SU(3)_u \times SU(3)_d$ for these LQ have been considered.

- To be consistent with asymptotic freedom of QCD, only a handful of flavor charges for the $S_1$ and $S_3$ are available.

- To find the constraint on the models, the MFV expansion is required to be self-consistent and perturbative, bounding the allowed couplings. Several models are thereby ruled out, and some others remain viable only within the context of a two-Higgs doublet model at large $\tan \beta$.

- Constraints on the MFV LQ models from $b \to s \nu \bar{\nu}$ and electroweak precision measurements were then considered. Some models were found to be ruled out by either of these constraints, or both.

As a result, all possible flavor quantum numbers for the scalar LQ are ruled out as an explanation of the $R_{D^{(*)}}$ anomaly within the assumption of Minimal Flavor Violation.


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91


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