Dark Z' Search and (Un)Natural Supersymmetric Models

by

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Abstract

The Standard Model (SM) of particle physics is the most successful theory of particles and fundamental laws of nature. It has been tested numerous times for over forty years. Yet, it is considered to be incomplete. It does not incorporate gravity and does not explain dark matter or dark energy. The parameters in the SM are ad-hoc in nature and the Higgs mass is unstable to quantum corrections. The SM covers only $\sim 5\%$ of the energy-matter content of the cosmos.

Dark matter (DM) has been indirectly observed via its gravitational effects on ordinary matter. Currently, there are no acceptable results from terrestrial experiments that can explain the particle properties of the DM. We consider several production models for a dark gauge boson, Z' that mediates a dark force. We find that by introducing new cuts, we can optimize dilepton resonance and MET searches at the Large Hadron Collider that can efficiently look for the Z' of mass of $\mathcal{O}(100\text{ GeV})$.

Supersymmetry (SUSY) is the most widely used framework for beyond the SM framework. It provides a good DM candidate and can achieve better gauge coupling unfication compared to the SM. For it to solve the hierarchy problem in a natural way, SUSY is expected to show itself at a scale of a few hundred GeVs. We present a model that combines two different SUSY breaking mechanisms allowing gaugino masses of $\mathcal{O}$(TeV) and yet preserving naturalness in the theory. For this purpose, we introduce messenger fields resulting in a compressed gaugino spectrum. This more compressed spectrum is less constrained by LHC searches and allows for lighter gluinos. In addition to the model, we present gaugino pole mass equations that differ from (and correct) the original literature.

We also consider the case where SUSY is not associated with the weak scale and solves the hierarchy problem by fine-tuning while retaining its other appealing features. A Mini-Split SUSY model is presented with SUSY scalars, m_{sc} in the mass range of $100-1000$ TeV. Higgsino masses, if not at the Planck scale, should generically appear at the same scale. The gaugino mass contributions from anomaly mediation, with the heavy Higgsino threshold, generally leads to a more compressed spectrum than standard anomaly mediation, while the presence of extra vector-like matter near m_{sc} typically leads to an even more compressed spectrum. Heavy Higgsinos improve gauge coupling unification relative to the MSSM. This model achieves the experimentally observed mass of Higgs and has a DM candidate.

Advisor: David E. Kaplan

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Chapter 1

Introduction

1.1 The Standard Model

The Standard Model (SM) of particle physics is a theoretical framework that was formed more than four decades ago. It is written in the language of quantum field theory and describes the fundamental interactions of the particles in the universe. It incorporates three fundamental forces of nature - electromagnetism, weak nuclear force and strong nuclear force.

All the particles that have been detected till today are described by the SM. The list of detected particles is comprised of spin-1 particles that mediate forces, spin-1/2 particles like quarks and leptons that make up the ordinary matter, and the spin-0 higgs that is instrumental in generating mass. Up, down, top, bottom, strange and charm are the quarks and the electron, muon, tau and the associated neutrinos make up the leptons. Gluons mediate the strong nuclear force between the quarks. W^+, W^- and Z mediate the weak force between the quarks and the leptons. The most well known of the three forces described by the SM is the electromagnetic force mediated by photons between quarks and charged leptons.

Even though, the SM has survived numerous experiments [2–4], it is not a complete description of the nature. In the many successes of the SM, the shortcomings stand out. Apart from fundamental questions like why do parameters like couplings and the number of particles have the value that they have in SM, the most glaring shortcoming is the absence of gravity in the SM framework. Thus, it can only be an effective theory up till the Planck scale ($\sim 10^{19} GeV$). SM also does not account for all the mass-energy content of the cosmos [5, 6]. Various cosmological studies imply that the ordinary baryonic matter is only 20% of the matter content of the universe. The rest of the matter is called the non-baryonic "dark matter" (DM) as it only weakly couples to ordinary matter and can't be "seen" in the ordinary sense of the word.

SM suffers from the "hierarchy problem". The Higgs vacuum expectation value (vev) is susceptible to quantum corrections that can drive it up to the Planck scale instead of the weak scale that is associated with SM. SM also does not account for neutrino oscillations. Astrophysical studies made in the late 1990s showed that we live in a cosmos that has an accelerating expansion [7]. The expansion seems to be fueled by a constant vacuum energy density called the "dark energy" and most popularly parametrised as the cosmological constant in Einstein's general theory of relativity.

SM has no explanation for the nature of dark energy that constitutes $\sim 70\%$ of the energy-matter content of the cosmos [8, 9].

Clearly, a lot of work needs to be done to improve upon the SM framework to solve the aforementioned problems. In this thesis we propose new search strategies for detecting a dark gauge boson in the DM sector and some phenomenological models in the SUSY framework that not only provide a DM candidate but come with other advantages that the SM lacks. In the following sections, we give a brief overview of these topics.

1.2 Dark Matter

A large number of astronomical and cosmological experiments/observations indicate that DM makes up to $\sim 25\%$ of the energy-matter content of the cosmos. Studies of the cosmic microwave background power spectrum [10], galaxy rotation curves, and other observations indirectly point to the presence of DM based on DM's gravitational interactions. These experiments do not comment on the mass or nature of the DM particles. Finding the value of such parameters is currently in the domain of experiments that are categorised under direct detection, indirect detection and detection at colliders. All these experiments assume weak interactions between DM and SM particles.

If the energy density of DM is a relic of freeze-out from the thermal bath of the early Universe, then the velocity-averaged cross-section for dark matter annihilation at freeze-out was ~ 1 pb. This is achieved if the annihilation process involves coupling constants $\lesssim .1$ and masses $\lesssim 100$ GeV, suggesting that the dark matter mass is around the weak scale, kinematically accessible to particle accelerators such as the Large Hadron Collider (LHC). Such particles are called the weakly interacting massive particles (WIMP). This "WIMP miracle" is a prime motivation for considering dark matter production at the LHC.

Direct and indirect detection experiments will only probe the stable DM particle. If the DM sector is minimal in nature, i.e., there is only one kind of DM particle, then all three types of experiments - direct, indirect and collider can try to access that particle. On the other hand, if the DM sector is non-minimal in nature, i.e., the DM sector is more complex in terms of types of dark matter particles, then only the collider experiments can access DM particles other than the final stable DM particle that builds the thermal relic density.

Current searches for dark matter at the LHC focus on pair production of invisible DM particles plus radiation from the initial state in the form of jets, photons, or electroweak bosons. The resulting "monojet", "monophoton", etc. signatures have considerable SM background, but still allow for constraints to be placed on dark matter interactions with quarks and gluons [11]. These can be compared to analogous bounds from elastic DM scattering in detectors (direct detection) and astrophysical DM annihilation (indirect detection).

As mentioned earlier, this program is appropriate for the minimal assumption of a single DM particle and no other new physics. However, when one goes beyond this minimal framework, other types of collider searches may provide much more powerful probes. A familiar example is that of supersymmetric models, in which searching for squarks and gluinos decaying to a neutralino DM candidate is usually a far more effective probe of the new physics than searches for direct neutralino production. More generally, any new particles associated with dark matter can provide additional collider signatures which may greatly enhance the prospects for discovery. Note that the same enhancement does not extend to direct and indirect detection of dark matter, which are only sensitive to the actual cosmological relics. In this respect colliders provide a unique window into the physics associated with dark matter.

In the associated chapter based on work in progress with Reinard Primulando and Prashant Saraswat, we propose new searches for a dark sector boson, Z' in the mass range of a few hundred GeV in the dilepton resonance + MET channel.

1.3 Supersymmetry

Supersymmetry (SUSY) is the most widely studied BSM physics paradigm. SUSY is a framework in which every particle with a half-integer spin has a SUSY partner with an integer spin and vice-versa. The supersymmetric partners of SM matter particles are called the "**s**particles". The supersymmetric partners of the gauge bosons are called "gauginos" and "higgsinos" are the partners of the higgs boson. Since no superpartners have been observed yet, they must be heavy to have evaded the existing experimental searches. This implies that SUSY is not an exact symmetry – particles and their superpartners have different masses, thus it has to be "broken". There are many viable ways in which this mass splitting is accomplished, each with its own advantages and drawbacks. A few SUSY breaking mechanisms will be commented on later.

1.3.1 Minimal Supersymmetric Standard Model

Most SUSY phenomenological models are based on the Minimal Supersymmetric Standard Model (MSSM). Weak scale supersymmetric extensions to the SM naturally cancel the quantum effects (upto some logarithmic dependence on the SUSY breaking mechanism) that enormously enhance m_H and stabilizes the Electroweak scale around $\mathcal{O}$(100 GeV). This is the prime motivation behind weak scale SUSY.

We require SUSY breaking to produce heavy SUSY particles that have eluded detection thus far. The SUSY breaking processes are parametrized by m_{soft}. The SUSY preserving sector helps in keeping $m_H \sim \mathcal{O}(100\text{ GeV})$ naturally. Since, the corrections to m_H depend on m_{soft}, we require m_{soft} to be about the same order as m_H to preserve the cancellations in the quantum effects that were provided by the SUSY preserving terms.

The SUSY preserving sector contains a parameter universally referred to as μ

which denotes the mass of the superpartner of the Higgs. This parameter plays an important role along with SUSY breaking parameters in determining the spectrum (particle content of the theory) of the theory. Since, all the discovered particles rely on m_H for their mass values, we have to make sure that $\mu \sim \mathcal{O}(m_{soft})$ with both of them being at most of the $\mathcal{O}(1 \text{ TeV})$ to produce $m_H \sim \mathcal{O}(100 \text{ GeV})$.

There is no clear reason as to why a SUSY preserving mass scale (μ) should be of the same order as a SUSY breaking mass scale (m_{soft}). There is no universally accepted mechanism of generating the μ term of the right mass scale. This puzzle is known as the "μ problem". One way of approaching this problem is to extend the simplest SUSY models by incoroporating mechanisms where μ directly depends on m_{soft}.

SUSY has other advantages. It proposes a natural dark matter candidate in the form of the lightest stable supersymmetric particle (LSP) and MSSM has better gauge coupling unification compared to SM.

To explain the dearth of SUSY signals at past and present colliders, SUSY has to be broken to give superpartner masses that evade current bounds. A successful SUSY breaking mechanism has to pass constraints related to flavor-changing neutral currents and should preserve gauge coupling unification. Two SUSY breaking mechanisms that satisfy these requirements are anomaly-mediated supersymmetry breaking (AMSB) and gauge-mediated supersymmetry breaking (GMSB), each with its own set of accompanying problems.

AMSB is a very appealing SUSY breaking mechanism since it is flavor blind, UV insensitive, and highly predictive [12, 13]. In the minimal AMSB setup, the ratio of the gluino mass to the LSP wino mass is about a factor of ten, meaning that this model is well covered by conventional SUSY searches that rely on hard jets and large missing transverse energy. The null results from such searches based on analyses of all of the 7 TeV data [14, 15] collected at the LHC and some or most of the 8 TeV data [16–19] have placed a lower bound on the gluino mass of about 1.5 TeV for many vanilla models, leading to increased tension with naturalness. Furthermore, AMSB models suffer from the tachyonic slepton problem and the μ problem.

GMSB is another widely studied SUSY breaking mechanism. SUSY breaking is communicated from a hidden sector to the superpartner fields via messenger fields that have interactions with SM gauge bosons. To preserve gauge coupling unification, the messengers should come in complete representations of $SU(5)$. GMSB models are predictive and do not suffer from flavor problems. However, they have a gravitino LSP which can be a problem in the context of cosmology. In addition, they also suffer from the μ problem.

In chapter 3 we present work based on [20], written in collaboration with David E. Kaplan and Tom Zorawski, which utilizes advantages of both AMSB and GMSB but also resolves the tachyonic slepton problem. We are able to get a compressed gaugino spectrum that is not as constrained as vanilla MSSM models.

1.3.2 High Scale SUSY

Weak scale SUSY is currently not supported by the experimental results. Further, with the discovery of a SM Higgs-like particle of mass of about 126 GeV implies superpartner masses of $\mathcal{O}(10TeV)$ [21], which is at odds with the idea of naturalness. Throwing away the naturalness idea but still keeping other advantages of SUSY like gauge coupling unification and a DM candidate is an approach taken in Split SUSY models [22–24].

At this juncture one wonders if the mass of the Higgs is an accidental tuning of parameters. Of course, completely natural supersymmetric theories may still turn out to describe physics at the TeV scale, and there have been no shortage of models of this sort proposed recently in response to null-results for new physics from the LHC. It is however fair to say that these models are rather elaborate. Many of these theories are actually just as fine-tuned as more conventional versions of supersymmetry, but the tuning is more hidden.

In the mini-split SUSY scenario, supersymmetric scalars have masses in the range of $\sim 100 - 1000$ TeV. Gaugino masses come out in the range of $\mathcal{O}(1 - 10 \text{ TeV})$ as a result of incorporating AMSB. Such a model is presented in detail in Chapter 4 based on [25], written in collaboration with Nima Arkani-Hamed, David E. Kaplan, Neal Weiner and Tom Zorawski.

Chapter 2

Probe for Dark Sector Dynamics at the LHC

2.1 Introduction

In this chapter we discuss the possible interactions of a dark sector boson leading to decays to SM leptons and missing energy. There is compelling evidence that most of the matter in the Universe is composed of nonbaryonic particles, the dark matter (DM), the nature of which is otherwise unknown.

In this work we specialize to the case where all new particles are gauge singlets under the Standard Model, forming a "dark sector." Although the model space for such a dark sector is vast, some well-motivated assumptions greatly narrow down the possible collider phenomenology. Consistent with renormalizable field theory, we can consider the new particles to be either fermions, scalars or gauge bosons. Since the coupling of these states to the SM is generally very weak, when dark sector particles are produced at colliders they will tend to cascade decay within the dark sector until that is no longer kinematically possible. In particular, particles with arbitrarily weak couplings to the SM can still be produced through decay of or radiation off of other dark sector particles. If the theory preserves baryon and lepton number, then the lightest dark fermion will be absolutely stable and appear as missing energy at colliders. Dark bosons however may not be protected by any quantum numbers and could decay into the Standard Model. Although this decay width may be small due to weak couplings, the branching ratio to the SM is of course unity if there are no kinematically allowed decay channels in the dark sector. This allows for visible collider signatures from extremely weakly coupled particles.

The collider signatures of a dark gauge boson, or Z', can be particularly striking. The most relevant possible interaction of a dark vector boson with the Standard Model is through the kinetic mixing "portal", i.e. the operator $B_{\mu\nu}X^{\mu\nu}$ where $B^{\mu\nu}$ is the SM hypercharge field strength and $X^{\mu\nu}$ is the dark gauge boson field strength. As a result of this mixing the Z' will couple to the SM hypercharge current and thereby decay into pairs of fermions, including an $O(1)$ branching ratio to leptons. This gives a distinctive and easily measured dilepton resonance in collider events. Direct $2 \to 1$ production of Z's has been searched for extensively at colliders, but the SM Drell-Yan background limits the sensitivity of such searches to weakly coupled Z's.

However, dark sector cascades will tend to produce Z''s in association with missing transverse momentum (MET) from invisible dark fermions. In this work we show that leveraging this large MET would allow for searches with very low background and high acceptance for a broad class of dark sector models.

In the next section we discuss the phenomenology of a kinetically mixed Z' including the collider signature and the bounds of the the model. In section III, the signal production at LHC and the possible production scenarios are discussed. Finally, we conclude with a discussion and comparision of the current bounds with the bounds obtained from the proposed cuts in this chapter.

2.2 A Kinetically Mixed Z'

We consider a dark sector where the fermion sector is charged under a $U(1)_D$. The dark gauge boson of the $U(1)_D$ mixes with the SM $U(1)_Y$ via kinetic mixing. The Lagrangian of the model is given by

$$\mathcal{L} = \mathcal{L}_{SM} - \frac{1}{4}\tilde{F}_{\mu\nu}\tilde{F}^{\mu\nu} + m_{Z'}^2\tilde{X}_\mu\tilde{X}^\mu + \frac{\epsilon}{2}\tilde{F}_{\mu\nu}\tilde{B}^{\mu\nu}, \tag{2.1}$$

where $\mathcal{L}_{SM}$ is the standard model Lagrangian, $\tilde{X}^\mu$ is the dark gauge boson, $\tilde{F}^{\mu\nu}$ is the field strength tensor of $\tilde{X}^\mu$, and $\tilde{B}^{\mu\nu}$ is the field strength tensor of $U(1)_Y$. $m_{Z'}$ is the mass of dark gauge boson which can arise from various mechanism, e.g. a spontaneous breaking of a dark Higgs or Stueckelberg mechanism.

In order to diagonalize the kinetic and mass term of the full Lagrangian, we redefine the gauge fields to be

$$\begin{aligned}
\tilde{A}_3^\mu &\rightarrow s_W A^\mu + c_\alpha c_W Z^\mu + s_\alpha c_W Z'^\mu, \\
\tilde{B}^\mu &\rightarrow c_W A^\mu - (c_\alpha s_W + \frac{\epsilon}{\sqrt{1-\epsilon^2}} s_\alpha) Z_\mu - (s_\alpha s_W - \frac{\epsilon}{\sqrt{1-\epsilon^2}} c_\alpha) Z'^\mu, \\
\tilde{X}^\mu &\rightarrow -s_\alpha Z^\mu + c_\alpha Z'^\mu,
\end{aligned} \tag{2.2}$$

where $\tilde{A}_3^\mu$ is the third component of the $SU(2)_L$ gauge boson; A^μ, Z^μ, Z'^μ are the mass eigenstates of photon, Z-boson and Z' respectively; s_W and c_W are the sine and cosine of weak mixing angle. The variables s_α and c_α are the mixing angle between the Z' and the SM Z-boson. In the limit of $\frac{\epsilon s_W}{1-m_{Z'}^2/m_Z^2} \ll 1$ they are given by

$$\begin{aligned}
s_\alpha &= \frac{\epsilon s_W}{1-m_{Z'}^2/m_Z^2} + \mathcal{O}(\frac{\epsilon s_W}{1-m_{Z'}^2/m_Z^2})^2 \\
c_\alpha &= \sqrt{1-s_\alpha^2},
\end{aligned} \tag{2.3}$$

where m_Z is the SM Z-boson mass.

For Z' mass between 1 GeV to 200 GeV, the electroweak precision measurements bound ϵ to be < 0.03 [26]. A bound from LHC Drell-Yan production of Z' was obtained in [27], which ϵ is bounded to be $\epsilon < 0.05$ for $m_{Z'} < 1$ TeV assuming that Z' does not decay to dark sector particles.

Z' decays to SM sector via the mixing with $U(1)_Y$ gauge boson. The decay width of Z' to SM fermion pair is given by:

$$\Gamma_f = \frac{1}{12\pi m_{Z'}} \sqrt{1-\frac{4m_f^2}{m_{Z'}^2}} \left((c_{fV}^2 + c_{fA}^2) m_{Z'}^2 + 2(c_{fV}^2 - 2c_{fA}^2) m_f^2 \right), \tag{2.4}$$

where m_f is the fermion mass and the coefficient c_{fV} and c_{fA} are given in table 2.1.

$c_{\nu V}$	$\frac{1}{4}(g_3 - g_Y)$	c_{uV}	$\frac{1}{4}(g_3 + \frac{5}{3}g_Y)$
$c_{\nu A}$	$-\frac{1}{4}(g_3 - g_Y)$	c_{uA}	$\frac{1}{4}(-g_3 + g_Y)$
$c_{\ell V}$	$-\frac{1}{4}(g_3 + 3g_Y)$	c_{dV}	$-\frac{1}{4}(g_3 + \frac{1}{3}g_Y)$
$c_{\ell A}$	$\frac{1}{4}(g_3 - g_Y)$	c_{dA}	$\frac{1}{4}(g_3 - g_Y)$

Table 2.1: c_{fV} and c_{fA} coefficients for various SM fermions. g_3 and g_Y are defined as $g_3 = g\, c_W\, s_\alpha$ and $g_Y = -g'\left(s_W s_\alpha - \frac{c_\alpha \epsilon}{\sqrt{1-\epsilon^2}}\right)$, where g and g' are the $SU(2)_L$ and $U(1)_Y$ gauge couplings respectively.

The decay width of Z' to a pair of W^+W^- Is given by

$$\Gamma_W = \frac{g^2 c_W^2 s_\alpha^2}{12\pi m_{Z'}} \sqrt{1 - \frac{4m_W^2}{m_{Z'}^2}} \left(\frac{m_{Z'}^6}{4m_W^4} + \frac{4m_{Z'}^4}{m_W^2} - 17m_{Z'}^2 - 12m_W^2 \right), \tag{2.5}$$

where g is the $\mathrm{SU}(2)_L$ gauge coupling and m_W is the W mass.

For the cases of Darkstrahlung and Cascade Decay (discussed in sections 2.3.1, 2.3.2, respectively), we assume that Z' does not decay to other dark sector particles. The branching fraction for Z' shown in Fig 2.1 is independent of ϵ. The branching fraction to leptons in the plot is assumed to be twice the branching fraction to electron-positron pair as both muons and electrons are very light compared to Z'. The plot shows appreciable branching fraction to the leptonic channel. Thus Z' decay to SM leptons is preferred as it is a cleaner channel compared to other SM decay channels (eg, decay to b-quarks). Further, the dilepton decay channel will show a resonance at the Z' mass. Moreover, the decay of Z' is prompt, unless the value of ϵ is extremely small as shown in Fig. 2.2. Prompt decays allow easy reconstruction of the event vertex.

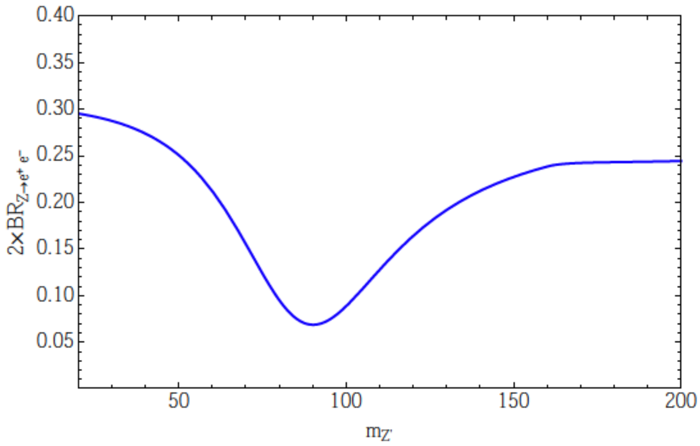

Figure 2.1: The branching fraction of the Z' to two leptons assuming Z' only decays to SM. $m_{Z'}$ is in GeV.

If $m_{Z'} > 2m_\chi$, Z' can also decays to dark matter pair. Assuming the dark matter to be a Majorana fermion, the partial decay width is given by

$$\Gamma_\chi = \frac{g_D^2}{24\pi m_{Z'}}\sqrt{1-\frac{4m_\chi^2}{m_{Z'}^2}}(m_{Z'}^2 - 4m_\chi^2), \tag{2.6}$$

where g_D is the coupling of Z' with dark sector fermions. For $g_D \sim 1$, Z' decays dominantly to dark matter. The branching fraction of Z' for some example benchmark points are given in Fig. 2.3. Increase in ϵ increases branching fraction to SM leptons whereas an increase in g_D decreases the branching fraction to SM leptons.

If the mass for dark matter χ and Z' is generated via a dark Higgs mechanism, the relevant Lagrangian is

$$\mathcal{L} \supset i\chi^\dagger\bar{\sigma}^\mu D_\mu\chi + |D_\mu\Phi|^2 + \frac{1}{2}y_\chi\Phi\chi\chi + h.c. + V(\Phi), \tag{2.7}$$

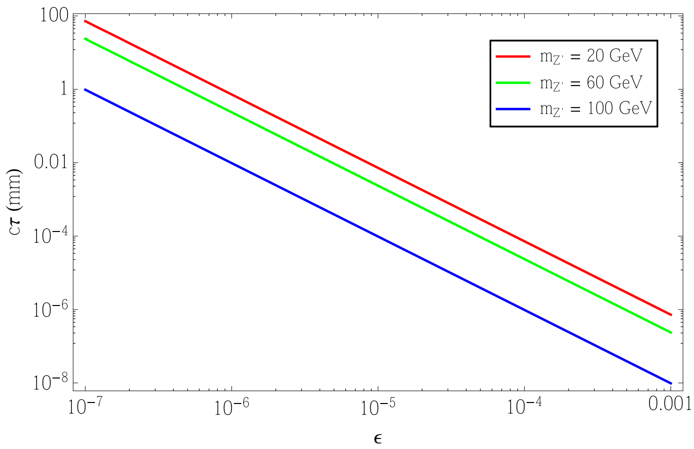

Figure 2.2: The decay length Z' as a function of kinetic mixing ϵ. We assume that the Z' decays 100% to SM. If the Z' decays to dark sector particles, the decay length becomes shorter.

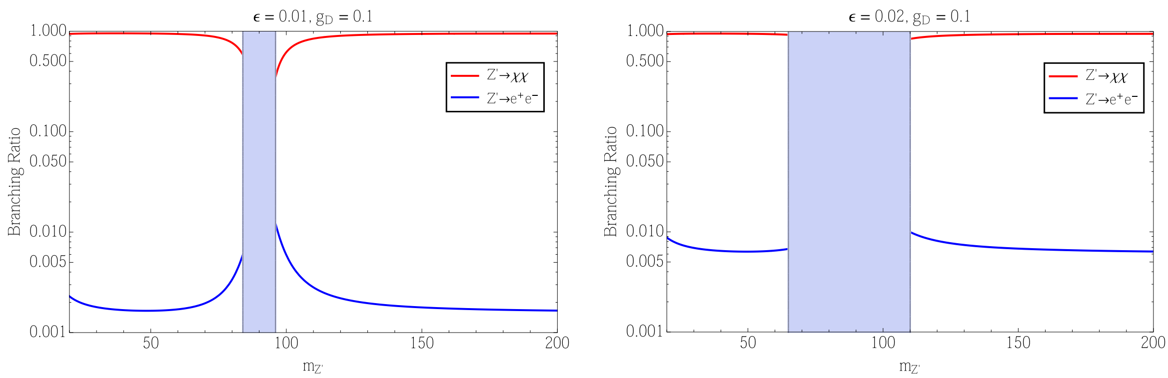

Figure 2.3: Branching fraction of Z' decaying to dark matter pair or a pair of electrons. The shaded area is excluded by precision electroweak measurements [26].

where Φ is the dark Higgs.

Here the dark matter is a Majorana fermion with a charge $q_\chi = -1$ under $\mathrm{U}(1)_D$ and the dark Higgs Φ has a charge $q_\Phi = -2$. After symmetry breaking, the dark Higgs acquire a vev v_D. Expanding Φ to be $\Phi = \frac{1}{\sqrt{2}}(v_D + \phi_D)$, we get the kinetic term for Φ to be

$$\begin{aligned} \mathcal{L} &\supset \frac{1}{2} q_\Phi^2 g_D^2 v_D^2 \tilde{X}_\mu \tilde{X}^\mu + q_\Phi^2 g_D^2 v_D h_D \tilde{X}_\mu X^\mu \\ &= \frac{1}{2} m_{Z'}^2 \tilde{X}_\mu \tilde{X}^\mu + q_\Phi g_D m_{Z'} h_D \tilde{X}_\mu X^\mu, \end{aligned} \tag{2.8}$$

where $m_{Z'} = q_\Phi g_D v_D$. The dark Yukawa coupling becomes

$$\begin{aligned} \mathcal{L} &\supset \frac{1}{2\sqrt{2}} y_\chi (v_D + h_D) \chi\chi + h.c. \\ &\supset q_\Phi g_D \frac{m_\chi}{m_{Z'}} h_D \chi\chi + h.c., \end{aligned} \tag{2.9}$$

where $m_\chi = \frac{y_\chi v_D}{\sqrt{2}}$ and we have substituted the expression for v_D to be $v_D = \frac{m_{Z'}}{q_\Phi g_D}$. The dark Higgs mixes with the SM Higgs and thus, can be produced at the LHC. The dark Higgs decay width to Z' and χ is given by

$$\Gamma(h_D \to Z'Z') = \frac{c_\alpha^4 c_\Phi^2 q_\Phi^2 g_D^2}{4\pi} \frac{m_{Z'}^2}{m_{h_D}^2} \sqrt{1 - \frac{4m_{Z'}^2}{m_{h_D}^2}} \left(3 + 2\frac{m_{h_D}^2}{m_{Z'}^2} - \frac{1}{2}\frac{m_{h_D}^4}{m_{Z'}^4}\right), \tag{2.10}$$

$$\Gamma(h_D \to \chi\chi) = \frac{c_\Phi^2 q_\Phi^2 g_D^2}{64\pi} m_{h_D} \frac{m_\chi^2}{m_{Z'}^2} \left(1 - \frac{4m_\chi^2}{m_{h_D}^2}\right)^{\frac{3}{2}}, \tag{2.11}$$

where c_Φ is the cosine of the mixing angle of SM Higgs and dark Higgs. In deriving the above expressions, we have assumed s_Φ and s_α to be small.

2.3 Z' Production at the LHC

Since the Z' couples weakly with the SM sector, the Drell-Yan production at the LHC is unobservable. The Drell-Yan production of Z' will be suppressed by $\mathcal{O}(\epsilon^2)$ compared to Drell-Yan production of SM Z-boson. On the other hand, the dark sector production cross section at the LHC is bounded to be at most $\mathcal{O}(10\text{ pb})$ from the monojet production cross section [11]. If the Z' is easily produced from the dark sector, for example from a radiation of Z' from dark sector particle, then the $Z'(\to \ell\ell)+\ \not{E}_T$ signature is present at a large number. This channel is considerably cleaner than the monojet; hence the production cross section bound can be better than monojet.

$Z'+$ jet is the dominant background at low MET but is negligible at MET above 40 GeV. The main background relevant to our analysis for this channel is W^+W^- where each W decays to a lepton and a neutrino. The leptons from the two W's has to be in the same flavor. The cross section times branching fraction to a same flavor lepton pair is given by 57.25 pb $\times 0.11^2 = 690$ fb. Since, the Z' is expected to have a narrow resonance, smaller bins in the m_{ll} analysis are recommended. Choosing the search mass window to be 5 GeV, the background is reduced by order of percent as shown in Fig. 2.4. Hence the bound on $Z'+\ \not{E}_T$ production cross section obtained from this search is $\mathcal{O}(1\text{ fb})$. As we discuss later, some additional cuts, e.g. $\not{E}_T$ cut can be introduced to eliminate most of the background while retaining a high signal acceptance.

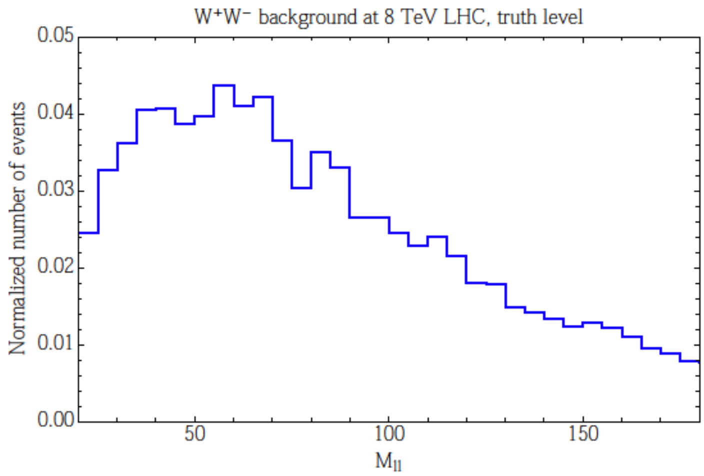

Figure 2.4: The distribution of dilepton invariant mass for W^+W^- background.

We study three differernt scenarios for Z' production - Darkstrahlung, Cascade decays and Dark Higgs. All of these prefer a mass of Z' of $\mathcal{O}(100$ GeV). We discuss the scenarios one by one

2.3.1 Darkstrahlung

In the darkstrahlung scenario, Z' is radiated from the dark matter, as shown in Fig. 2.5. The dark matter interaction with SM particle is described by an effective field theory (EFT) description

$$\mathcal{L} \supset \frac{\bar{q}\gamma^\mu\gamma^5 q\,\bar{\chi}\gamma_\mu\gamma^5\chi}{\Lambda^2}, \tag{2.12}$$

where q is the SM quarks and χ is the dark matter. This EFT captures scenarios where the dark sector particles are produced with significant boost and radiates Z'. The radiated Z' also receives some boost and thus, has some significant p_T.

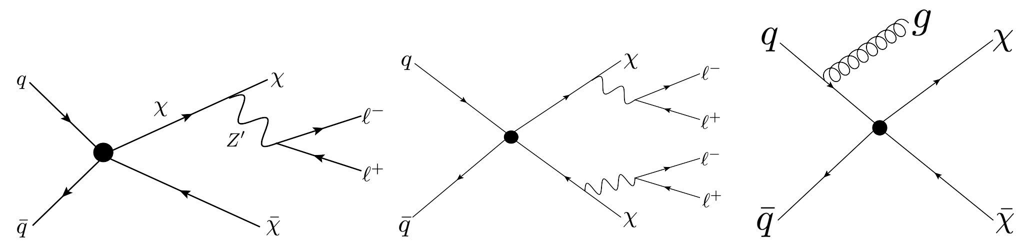

Figure 2.5: Feynman diagrams for $\ell^+\ell^-+\not{E}_T$, four leptons and monojet final states in the context of darkstrahlung scenario.

Fig. 2.6 shows missing energy distributions for the signal and the W^+W^- background. The background $\not{E}_T$ peaks around 40 GeV and becomes small for $\not{E}_T \gtrsim m_W$. On the other hand, the signal $\not{E}_T$ distribution extends beyond m_W, since Z' is boosted. A cut on high enough MET reduces the background to be essentially zero while retaining a high signal acceptance.

There are two other search channels that can also provide some additional bounds for this scenario. Monojet search, where the jet is radiated from the incoming quarks (see Fig. 2.5), can be a dominant bound when the coupling of dark matter and Z' is reduced. The current bound on the direct production of the dark sector at the LHC comes from the monojet + $\not{E}_T$ search [28]. This bound, when expressed in the terms of direct dark sector production, gives a $pp \to \chi\chi$ cross section of $\mathcal{O}(10\text{ pb})$ at 8 TeV LHC. Since the dark sector is charged under $U(1)_D$, the production cross section of $pp \to Z'\chi\chi$ is in order of $\mathcal{O}(\alpha_D/(2\pi)^3 \times 10\text{ pb} \times \text{BR}_{Z'\to\ell\ell}) \sim \mathcal{O}(\alpha_D \times 10\text{ fb})$.

On the other hand, if the DM-Z' coupling is large, an additional Z's can be radiated from DM leading to a four leptons in the final state, as shown in Fig. 2.5. In this case, multi lepton search can give an additional bound. The cross section for this channel can be approximated to be $\mathcal{O}(\alpha_D^2/(2\pi)^6 \times 10 \text{ pb} \times \text{BR}^2_{Z' \to \ell\ell}) \sim \mathcal{O}(\alpha_D^2 \times 0.01 \text{ fb})$. The number of events for this channel 20 fb^{-1} of luminosity is $\sim 0.1 \times \alpha_D$. Hence the multi lepton signature of this model is not going to be observed at the 20 fb^{-1} LHC8, unless α_D is beyond pertubativity. We will discuss the dependence of the bounds on α_D in more detail in the next section.

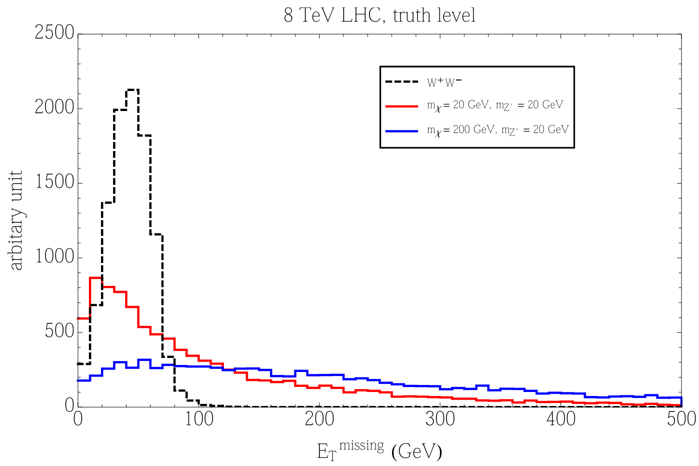

Figure 2.6: The MET distribution for signal and W^+W^- for a mass window $m_{ll} = 20 \pm 2.5$ GeV for darkstrahlung scenario.

2.3.2 Cascade Decays

This scenario mimics the previous scenario except Z' comes from a decay of $\psi \to \chi + Z'$, where ψ is another fermion in the dark sector, as shown in Fig. 2.7. The

Lagrangian in this scenario is given by

$$\mathcal{L} \supset \frac{\bar{q}\gamma^\mu\gamma^5 q\, \bar{\psi}\gamma_\mu\gamma^5\chi}{\Lambda^2} + \bar{\psi}\gamma^\mu\gamma^5\chi Z'_\mu + \text{ h.c..} \tag{2.13}$$

If there is no flavor diagonal dark fermions coupling to quarks and Z', neither monojet nor multi leptons channels are present. Hence only $\ell\ell+\not{E}_T$ channel can provide some bounds for this scenario.

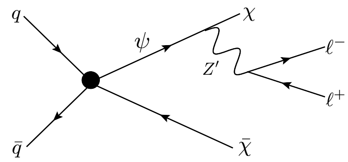

Figure 2.7: Feynman diagrams for $\ell^+\ell^-+\not{E}_T$ final state in the context of cascade decay scenario.

Like in the darkstrahlung case, both χ and ψ get a significant amount of boost. The events tends to have large missing energy as shown in Fig. 2.8. Hence a high $\not{E}_T$ cut does not reduce the signal acceptance significantly.

2.3.3 Dark Higgs

Other effective model that we consider is a scalar, ϕ, decaying to two Z', where the Z' has a large invisible branching fraction. This model contains a new scalar charged under $\mathrm{U}(1)_D$ that mixes with the SM Higgs. We assume a mass hierarchy $m_\phi > 2m_{Z'} > 4m_\chi$ so that the scalar decays to a Z' pair followed by Z' mainly decaying to dark matter. Besides decaying to dark matter, Z' can also decay to SM fermion pair

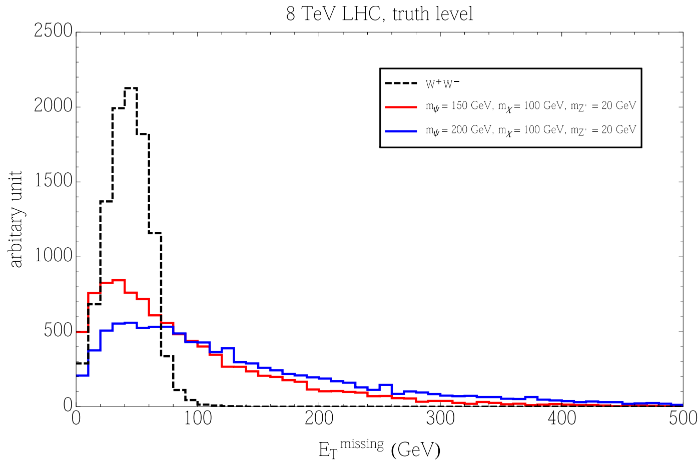

Figure 2.8: The MET distribution for signal and W^+W^- for a mass window $m_{ll} = 20 \pm 2.5$ GeV for cascade decay scenario.

via kinetic mixing. We are interested in a decay chain $\phi \to Z'Z' \to \chi\chi\ell^+\ell^-$ as shown in Fig. 2.9.

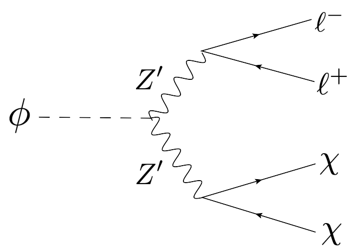

Figure 2.9: Feynman diagrams for $\ell^+\ell^-+ \not{E}_T$ final state in the context of dark Higgs scenario.

If the scalar mass is large and its mass is well separated from the Z' mass, the events will have a large missing energy. As mentioned in the previous scenario, the background MET distribution vanishes for $\not{E}_T \gtrsim M_W$ as shown in Fig. 2.10. Hence if

$\frac{m_\phi}{2} - m_{Z'} \gtrsim m_W$, we can put a $\not{E}_T$ cut where the signal region has a zero background. For a smaller mass splitting, the $\not{E}_T$ cut has to be lowered.

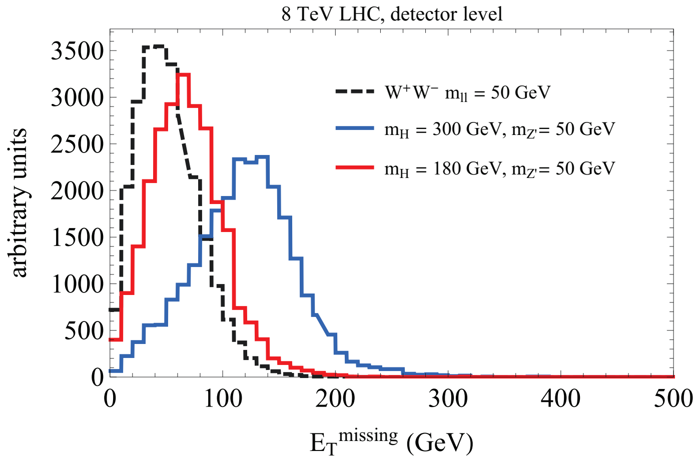

Figure 2.10: The MET distribution for signal and W^+W^- for a mass window $m_{ll} = 50 \pm 2.5$ GeV for dark higgs scenario.

Both monojet and multi lepton search can also provides additional bounds for this scenario. The strength of the bounds depends on the branching fraction of the Z' to missing energy and two leptons.

2.4 Proposed Search

2.4.1 Current Bounds

Both ATLAS and CMS collaborations have done some analysis considering final states of a lepton pair + $\not{E}_T$. CMS collaboration analyzed 5.3fb^{-1} of 8 TeV data to measure the W^+W^- cross section [29]. In the context of electroweakino productions, both of CMS [30] and ATLAS [1,31] have released their 8 TeV with ~ 20 fb^{-1} of luminosity. One of the ATLAS analysis [1] has released same flavor lepton pair invariant mass distribution which can be recast to obtain current bounds on Z' production.

The cuts employed by CMS [29] and ATLAS [1,31] analysis are given in table 2.2. The missing energy variable, $\not{E}_{T rel}$, is defined as

$$\not{E}_{T rel} = \begin{cases} \not{E}_T & \text{if } \Delta\phi_\ell > \pi/2 \\ \not{E}_T Sin\Delta\phi_\ell & \text{if } \Delta\phi_\ell < \pi/2 \end{cases}, \tag{2.14}$$

where $\Delta\phi_\ell$ is the azimuthal angle between the $\not{E}_T$ with the nearest lepton or jets.

We generated samples of WW, ZV and $t\bar{t}$ backgrounds as well as the signal at the parton level using Madgraph5 [32] followed by parton shower and hadronization simulation using Pythia 6 [33]. Additionally, pileup contributions are added to the sample with the average number of pileups $\langle\mu\rangle = 20.7$. We used Delphes 3 [34] to simulate the detector effect. The distributions of WW background sample for events satisfying ATLAS cuts are given in Fig. 2.11. We found that adding pileup to the simulated sample is necessary to reproduce ATLAS results. The agreements between ATLAS and our simulations is around 30%.

The ATLAS electroweakino search [1] shows the observed $m_{\ell\ell'}$ distribution of two same flavor leptons which is consistent with the SM prediction. In the plot, the $m_{\ell\ell'}$ distribution is divided into bins with 10 GeV width. While the width of the bin is not optimized for the Z' search, bounds on Z' production can still be obtained.

	ATLAS SR WWa	CMS
$p_{T,e}$	> 10 GeV	> 10 GeV
$\lvert\eta_e\rvert$	< 2.47	< 2.5
$p_{T,\mu}$	> 10 GeV	> 10 GeV
$\lvert\eta_\mu\rvert$	< 2.4	< 2.4
$p_{T,\text{leading lepton}}$	> 35 GeV	> 20 GeV
$p_{T,\text{second lepton}}$	> 20 GeV	> 20 GeV
$m_{\ell\ell'}$	> 20 GeV	> 12 GeV
$\lvert m_{\ell\ell'} - m_Z\rvert$	> 10 GeV	> 15 GeV for same flavor
$p_{T,\ell\ell'}$	> 80 GeV	> 45 GeV
$\not{E}_{T\,rel}$	> 80 GeV	> 45 GeV
jet veto	events with $p_j^T > 20$ GeV	events with $p_j^T > 30$ GeV
	and $\lvert\eta_j\rvert < 2.4$.	and $\lvert\eta_j\rvert < 4.7$.
	events with $p_j^T > 30$ GeV	events with $p_j^T > 15$ GeV,
	and $2.4 < \lvert\eta_j\rvert < 4.5$	$\lvert\eta_j\rvert < 4.7$ and $\Delta\phi_{\ell\ell',j} > 165^0$

Table 2.2: The cuts employed by ATLAS and CMS analysis. Both analysis require exactly two opposite sign leptons. In the table above m_Z is the mass of SM Z-boson and $\Delta\phi_{\ell\ell',j}$ is the azimuthal angle between the jet momentum and the dilepton system total momentum. $\not{E}_{T\,rel}$ is defined in the text.

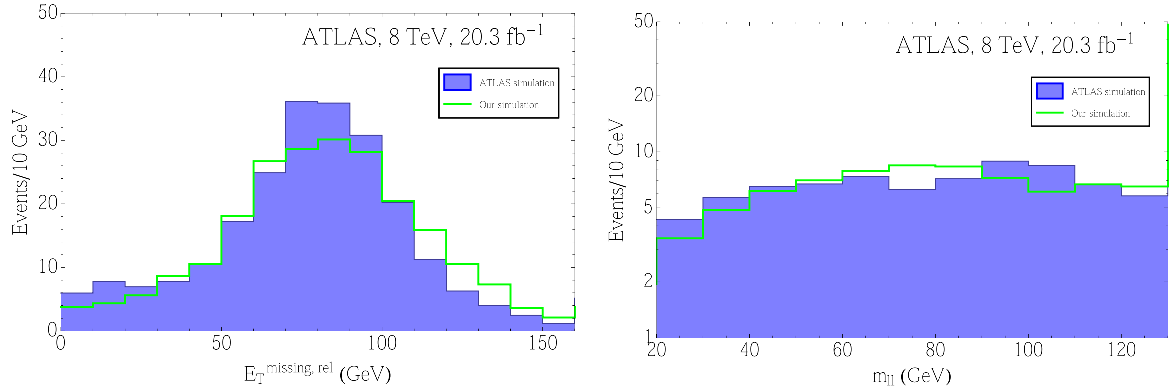

Figure 2.11: $\not\!\!E_{T rel}$ and $m_{\ell\ell}$ distributions for WW background. For the $\not\!\!E_{T rel}$ distribution, the ATLAS SRWWa cuts are applied except $\not\!\!E_{T rel}$ and $m_{\ell\ell}$. While for the $m_{\ell\ell}$ distributions, the cut on $m_{\ell\ell}$ is not applied.

The observed data and the analysis reported in [1] were used to genrate the two-sigma exclusion limits on $\sigma_{pp\to\ell^+\ell^-\chi\chi} \times BR(\text{fb})$ for the Darkstrahlung model as shown in Fig. 2.12. The limits are of $\mathcal{O}(few\text{fb})$ for $m_{Z'}$ ranging from $30 - 130$ GeV. The exclusion limits should stay similar for the cascade decay model since the only difference between the two models is that Z' is radiated from a dark matter fermion in the first model and a decay product in the latter model. In both models, we expect Z' to be boosted. We also produced the two-sigma exclusion limits on $\sigma_{pp\to\ell^+\ell^-\chi\chi} \times BR(\text{fb})$ for the Dark Higgs model as shown in Fig. 2.13. The limits are higher in this case. There is a region in which $m_{h_D} \sim 2m_{Z'}$ where Z' is not boosted enough resulting in low MET - not enough to pass the MET cut imposed in ATLAS analysis.

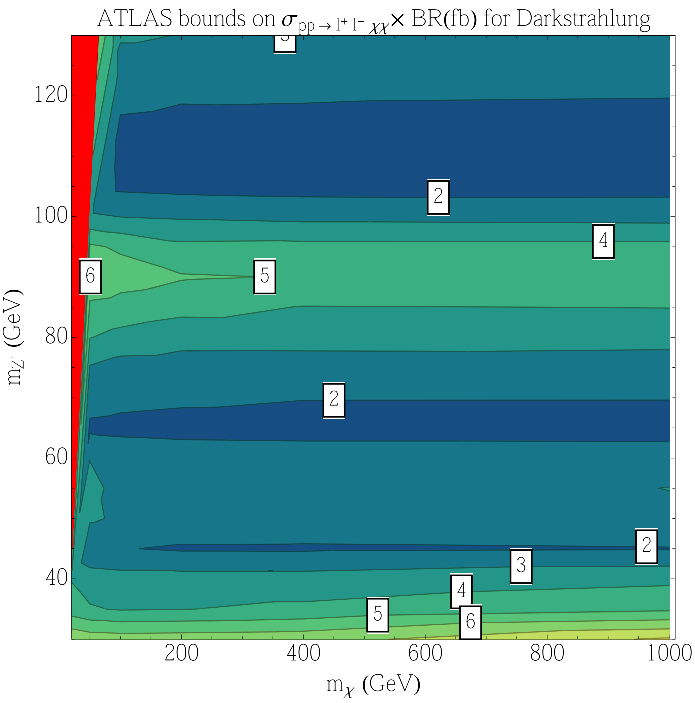

Figure 2.12: The bound obtained from recasting ATLAS electroweakino search [1] for darkstrahlung screnario. The red shaded area is the region which $m_{Z'} > 2m_\chi$ so that the Z' decays invisibly.

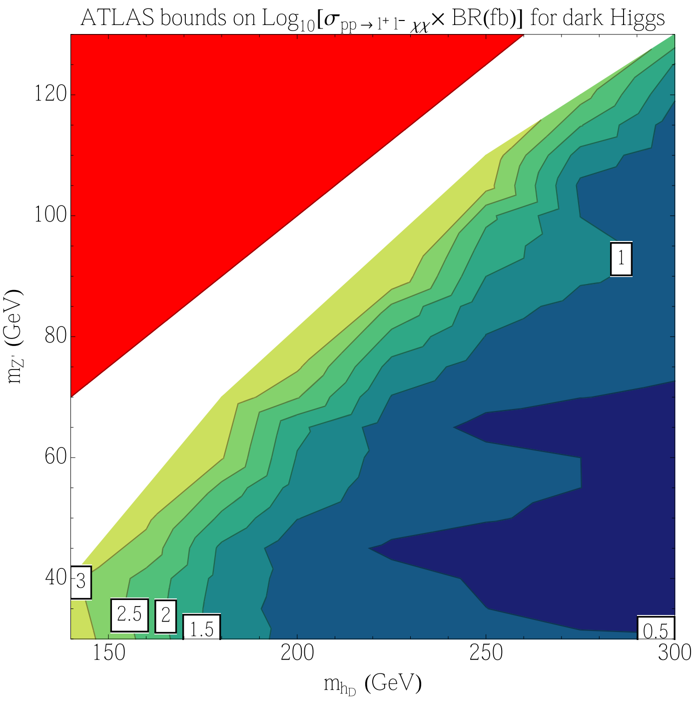

Figure 2.13: The bound obtained from recasting ATLAS electroweakino search [1] for the dark higgs screnario. The red shaded area is the region which $m_{Z'} > m_{h_D}/2$ preventing decay to Z'.

2.4.2 Optimized Bounds

The bound obtained in the previous subsection can be optimized by introducing new cuts. The improvements on the analysis are made in the following ways:

- A narrower $m_{\ell\ell}$ window,
- A harder $\not{E}_{T\,rel}$ cut,

Since the Z' is a narrow resonance, the width of $m_{\ell\ell}$ distribution depends on the detector resolution. A signal simulation of Higgs decay of two muons estimates a FHWM of 5 GeV [35]. The resolution can change as a function of $m_{\ell\ell}$. However since we are interested in Z' in a mass range of $10 - 100$ GeV, we use a 5 GeV window for defining the signal region.

As discussed in subsection 2.3.1 and 2.3.2, in the Darkstrahlung and cascade decay scenarios, the Z' is boosted and the events generally have a significant missing energy. On the contrary, background MET distribution becomes small for $\not{E}_T \gtrsim 80$ GeV. This suggest that the search can be optimized by introducing a substantial $\not{E}_T$ cut to reduce the background to be essentially zero.

We introduce a MET cut at 100 GeV for the Darkstrahlung model resulting in a reduction in the background but still maintaining enough efficiency for the signal. We simulated the SM background for the new cuts and produce the expected two-sigma exclusion limits on $\sigma_{pp\to\ell^+\ell^-\chi\chi} \times BR(\text{fb})$ for the Darkstrahlung model as shown in Fig. 2.14.

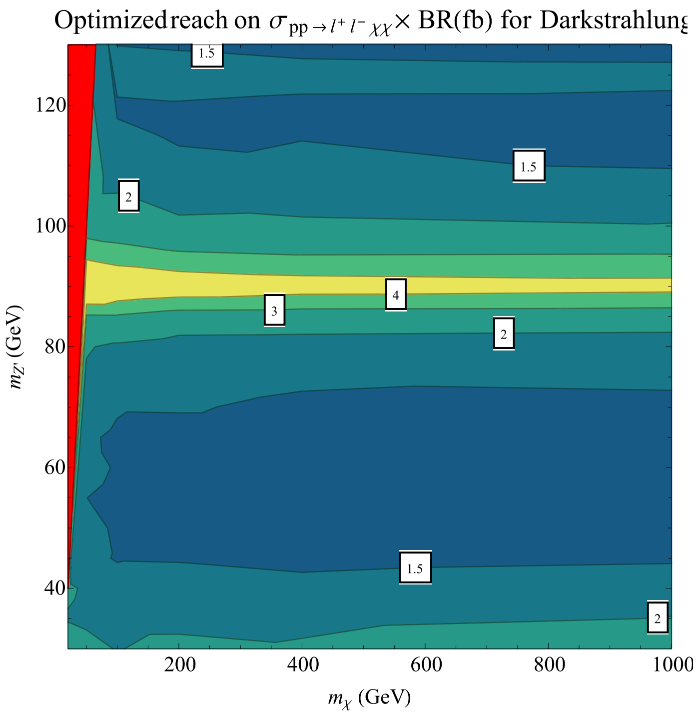

Figure 2.14: The expected bound obtained from optimizing cuts for the darkstrahlung screnario. The red shaded area is the region which $m_{Z'} > 2m_\chi$ so that the Z' decays invisibly.

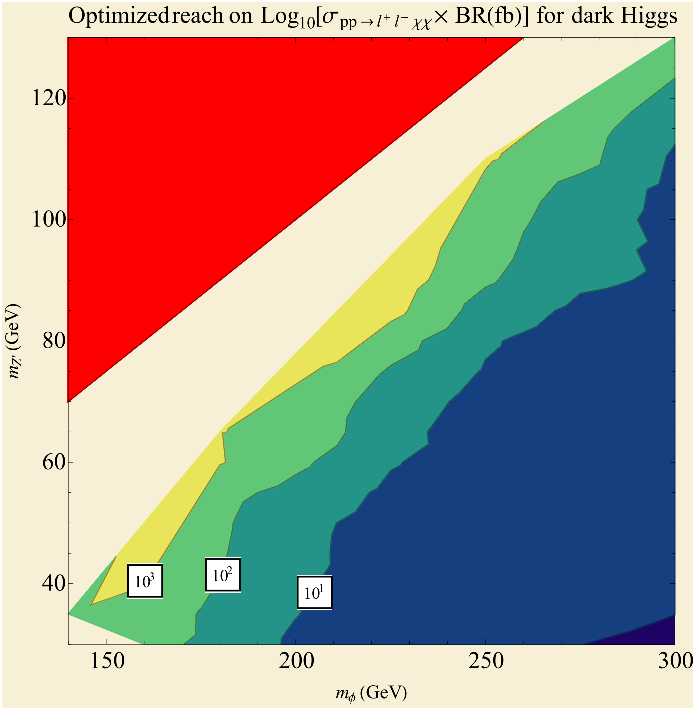

Figure 2.15: The expected bound obtained from optimizing cuts (High MET) for the dark higgs screnario. The red shaded area is the region which $m_{Z'} > m_{h_D}/2$ preventing decay to Z'.

For the Dark Higgs scenario too, we can apply a smaller m_{ll} window. For a large mass splitting between dark higgs and the Z', a high MET cut is desirable. On the other hand, in the case of the small mass splitting, the MET cut has to be lowered to probe the less boosted Z'. To solve this dilemma, we define two separate search regions - one with a MET cut of $p_T > 100$ GeV and the other with a low MET cut of $p_T > 50$ GeV. Using the simulated the SM background for the new cuts, the expected two-sigma exclusion limits on $\sigma_{pp\to\ell^+\ell^-\chi\chi} \times BR(\text{fb})$ for the Dark Higgs model are produced for the two different cuts. The limits for the high MET cut search are shown in Fig. 2.15 and the limits for the low MET cut search are shown in Fig. 2.16

2.4.3 Comparision between Monojet, Dilepton and Multi-lepton Searches

As mentioned earlier, the monojet and multi lepton search also bound dark-strahlung model. To compare the bound from difference searches, we pick a benchmark point $m_\chi = 100$ GeV and $m_{Z'} = 50$ GeV. Bounds on the effective mass scale and the production cross section are shown in Fig. 2.17.

The two plots in Fig. 2.17 are related by the inverse relationship between the production cross-section and Λ^4. Monojet production doesn't depend on the gauge coupling in the dark sector. The other two channels depend on the coupling but the dilepton channel dominates over the multiple lepton channel till the coupling increases

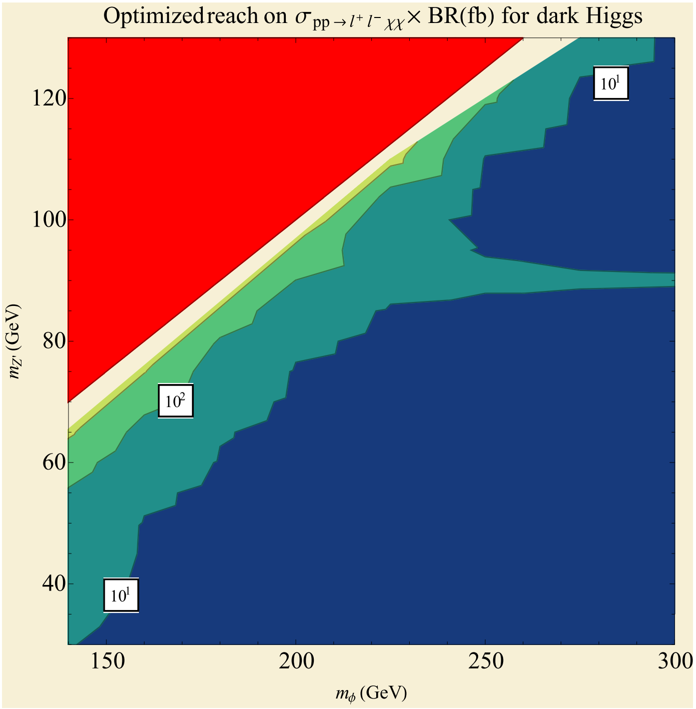

Figure 2.16: The expected bound obtained from optimizing cuts (Low MET) for the dark higgs screnario. The red shaded area is the region which $m_{Z'} > m_{h_D}/2$ preventing decay to Z'.

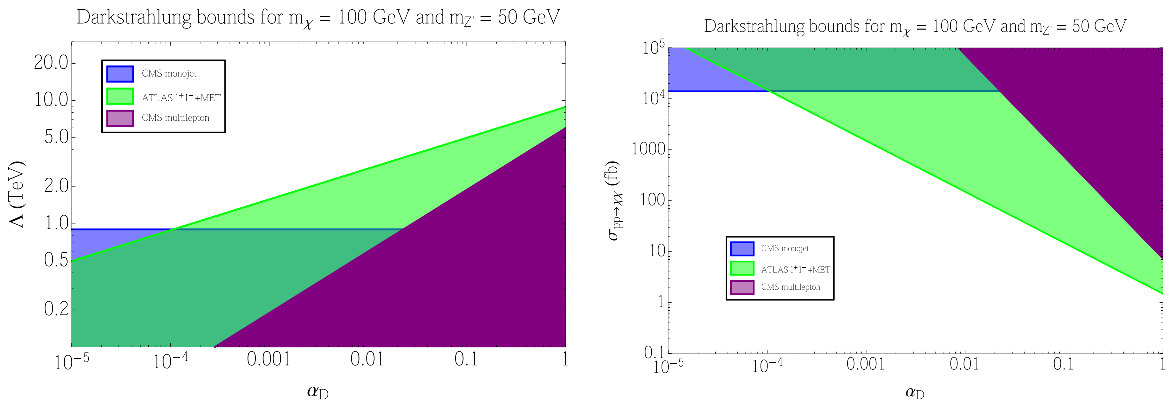

Figure 2.17: Bounds on Λ and cross-section from Z'+ MET search, monojet and multi lepton searches. The shaded area is excluded.

beyond perturbativity.

2.5 Conclusions

The continuous absence of DM detection in the direct detection and indirect detection experiments motivates the search for a non-minimal dark sector. Without being concerned with the exact model building details, we discussed the phenomenology of Z' production via radiating off a DM fermion and as a decay from a DM fermion or a dark higgs - scenarios that are being refererred to as Darkstrahlung, Cascade Decay and Dark Higgs. We are considering Z' decay to SM dileptonic channe plus METl that is cleaner compared to other decay channels and expect Z' to be a narrow resonance.

In the mass range of 20-120 GeV, Z' will be produced with high associated MET allowing to impose high MET cuts resulting in better exclusion limits on cross-section times branching ratio for dileptonic decay of Z'. Using the data from the current ATLAS electroweakino search, a bound of a few pb can be placed.

With a MET cut of $p_T > 100$ GeV in case of Darkstrahlung, the bounds could be improved by a factor of 3-4 times the current ATLAS bounds. Similar results should be obtainable in case of Cascade Decays as the phenomenology of decays is similar. In case of Dark Higgs scenario, two separate search regions are created that are complimentary. The search regions correspond to MET cuts of $p_T > 100$ GeV and $p_T > 50$ GeV. With these optimised cuts, the bounds could be improved by some orders of magnitude.

In case of Darkstrahlung, the dilepton channel associated with MET performs better at placing bounds compared with monojet and multiple lepton searches. We show that dilepton resonance + MET is a generic and powerful signature to probe a non-minimal dark sector.

Chapter 3

Gaugomaly Mediation Revisited

3.1 Introduction

In this chapter, we present a model that combines gauge and anomaly mediation to solve the tachyonic slepton problem and that also relaxes the constraints on gaugino masses by compressing the spectrum. This hybrid approach is not new. Refs. [36,37] first showed that the D-type gauge mediation of Poppitz-Trivedi [38] could be simply combined with AMSB to solve the tachyonic slepton problem, although they did not specify the origin of the messenger masses. Ref. [39] studied this further and gave it the name 'gaugomaly' mediation. A less direct solution to the slepton problem was proposed in ref. [40], which developed extended anomaly mediation (EAM) by arranging for the messengers to get masses directly from anomaly mediation through Giudice-Masiero (GM) type terms [41]. This deflects the gaugino masses off of the

AMSB trajectory, changing the scalar masses through running. The EAM setup was itself extended in [42, 43] with the addition of a singlet to yield realistic spectra.

We take the approach that such singlets are unnatural, so we are led to consider EAM with the D-type GMSB of gaugomaly mediation, which surprisingly had not been explored previously. In addition, we investigate the effect of the messengers in compressing the gaugino spectrum, an interesting aspect not discussed in the references above that deserves attention on its own. Integrating out the messengers takes the gaugino masses off the AMSB trajectory and gives threshold corrections that modify the masses at leading order, with the ratios sensitive to the number of messenger pairs. The result is a compressed gaugino spectrum with the mass splitting between the gluino and the LSP depending on the number of messenger pairs. The limits on the allowed gaugino masses are significantly weakened due to the squeezing of the spectrum [44]. This framework can also bring models that are otherwise beyond the reach of the LHC to within its reach in certain cases as the gluino becomes only 1-2 times heavier than the LSP (the wino or the bino depending on the number of messengers). In addition to the model, we present gaugino pole mass equations that differ from (and correct) the original literature.

This chapter is structured as follows. First, we briefly review the minimal AMSB framework. Although this has already been discussed in [45], we present the details here because our equations differ slightly. We introduce messengers in Sec. 3.4 in the context of D-type gauge mediation as a solution to the tachyonic slepton problem. In

Sec. 3.5, we discuss the gaugino spectra and related phenomenology independently from the scalars, focusing on the compression resulting from the messenger threshold corrections to the gaugino masses. Then, in Sec. 3.6, assuming that the μ problem has been solved, we give complete example spectra for our model with the scalars. We conclude in Sec. 3.7 by presenting a simple approach to addressing the μ problem in this framework and briefly discuss how it can be improved by incorporating some ideas found in the literature.

3.2 Review of AMSB

In the pure AMSB case, the gaugino masses are generated at the one loop level and the scalar masses squared are generated at two loops. The equations for these soft masses are the solutions to the RGEs for the gauginos and sfermions and thus are valid at all energy scales. The running gaugino masses are given by

$$m_i = \frac{\beta(g_i)}{g_i} m_{3/2} \tag{3.1}$$

This formula is valid at all scales, meaning that if we wish to calculate the masses in the IR we only need to know the values of the couplings at that scale. Using the one-loop beta functions, we obtain

$$m_1 = 11\frac{\alpha_1}{4\pi}m_{3/2} \quad (3.2)$$

$$m_2 = \frac{\alpha_2}{4\pi}m_{3/2} \quad (3.3)$$

$$m_3 = -3\frac{\alpha_3}{4\pi}m_{3/2}. \quad (3.4)$$

Note that throughout this paper we use the non-GUT normalization for hypercharge. Plugging in the weak scale values of the α_i, we find that the mass ratios $m_1 : m_2 : |m_3|$ are approximately 3.3 : 1 : 10. We assume here that μ is larger than any of the gaugino masses, so that the wino is the LSP in this minimal case. This large splitting between the gluino and the wino means that energetic jets will be produced in the cascade decay (we discuss phenomenology in a later section), leading to tight constraints on the masses. However, as pointed out in [**?**], quantum corrections here are significant, so we consider the pole masses computed at NLO. We discuss the calculation of the pole masses for the gauginos in the subsequent section.

The scalar masses in AMSB are given by

$$m_i^2 = -\frac{1}{4}\left(\frac{d\gamma_i}{dg_j}\beta_{g_j} + \frac{d\gamma_i}{dy_j}\beta_{y_j}\right)m_{3/2}^2, \quad (3.5)$$

where γ is the corresponding anomalous dimension, and β_g, β_y are the gauge coupling and yukawa coupling beta functions, respectively. Since the sleptons are charged only under non-asymptotically free gauge groups, they will have negative squared masses. This is the well-known tachyonic slepton problem in AMSB.

3.3 Gaugino Pole Masses

The full NLO expression includes contributions coming from α_3 and y_t in the two-loop beta functions as well as self-energy corrections. For the one-loop self-energies, our analysis follows closely the steps presented in [46].

3.3.1 Gauge Loops

The contribution to the self-energy from gauge boson loops is given by

$$\left(\frac{\Delta m_i}{m_i}\right)_{gauge} = \frac{\alpha_i}{4\pi} C(G_i) \left[4B_0(m_i, m_\psi, m_\phi) - 2B_1(m_i, m_\psi, m_\phi)\right], \tag{3.6}$$

where the Veltman-Passarino functions (defined as in [46]) are

$$B_0(m_i, m_\psi, m_\phi) = -\int_0^1 dx \ln \frac{(1-x)m_\psi^2 + x m_\phi^2 - x(1-x)m_i^2}{Q^2} \tag{3.7}$$

$$B_1(m_i, m_\psi, m_\phi) = -\int_0^1 dx\, x \ln \frac{(1-x)m_\psi^2 + x m_\phi^2 - x(1-x)m_i^2}{Q^2}, \tag{3.8}$$

with m_ψ the mass of the fermion in the loop, Q is the renormalization scale, and m_ϕ is the mass of the boson in the loop (vector or scalar). Throughout most of this paper, except where noted otherwise, we work in the $\overline{DR}$ scheme. For the gluino, $m_\phi = 0$ (the bosons in the loop are massless gluons), and we can evaluate the integrals directly, yielding

$$\left(\frac{\Delta m_3}{m_3}\right)_{gauge} = \frac{\alpha_3}{4\pi}C(G_3)\left(5 + 3\ln\frac{Q^2}{m_3^2}\right). \tag{3.9}$$

For the wino, we must use the full B functions, since m_W is of order m_2. The amount of diagrams differs for the neutral and charged wino (in fact, this will be the main source of the splitting between the lightest chargino and the LSP as described later). Here we use the empirical fit of [46] for the neutralino result.

3.3.2 Matter Loops

Next, we consider the contributions of the matter content to the gaugino masses. We assume that all squarks are degenerate with mass $m_{\tilde{q}}$, and that all sleptons are degenerate with mass $m_{\tilde{l}}$, with $m_{\tilde{q}}, m_{\tilde{l}} > m_3$ (this will be important when discussing the phenomenology). In addition, we approximate the fermion masses as zero. For the gluino, only quark/squark loops contribute, giving

$$\begin{aligned}
\left(\frac{\Delta m_3}{m_3}\right)_{matter} &= -12\frac{\alpha_3}{4\pi}B_1(m_3, 0, m_{\tilde{q}}), && (3.10)\\
-B_1(m_3, 0, m_{\tilde{q}}) &= -\frac{1}{2}\ln\frac{Q^2}{m_3^2} + I, && (3.11)\\
I &= \int_0^1 dx\, x\ln[rx - x(1-x)], \quad r = \left(\frac{m_{\tilde{q}}}{m_3}\right)^2. && (3.12)
\end{aligned}$$

It is straightforward to evaluate I

$$I = \frac{1}{2}\left(-2 + r + (r-1)^2 \ln|1-r| - (r-2) r \ln r\right). \tag{3.13}$$

Since the top mass is fairly large, we should examine whether neglecting it is a good approximation. Including the top mass gives an extra contribution $\frac{\alpha_3}{4\pi}\frac{m_t}{m_3} f(m_{\tilde{t}_1}, m_{\tilde{t}_2})$, where f is a function of the stop masses that is ~ 1. Since in our case $m_3 \sim 1$ TeV, this term is indeed suppressed.

For the wino and bino, we divide the matter contributions into fermion/sfermion, chargino/charged Higgs, and neutralino/neutral Higgs pieces. The sfermion piece is

$$\left(\frac{\Delta m_i}{m_i}\right)_{sfermion} = -2\frac{\alpha_i}{4\pi} S(R_\Phi) B_1(m_i, 0, m_\phi), \tag{3.14}$$

summing over all chiral supermultiplets $\Phi = (\phi, \psi)$ that couple to the wino/bino. To simplify B_1, we can take the wino/bino mass to be approximately zero. At first sight this doesn't seem to be valid, since the gaugino and sfermion masses are both of roughly the same order. However, since B_1 is multiplied by α_1 or α_2, which are both very small, any error will become negligible. We have in fact checked that this is so numerically. With this approximation we find

$$-B_1(0, 0, m_\phi) = -\frac{1}{2}\left(\ln\frac{Q^2}{m_i^2} - \ln\frac{m_\phi^2}{m_i^2} + \frac{1}{2}\right). \tag{3.15}$$

The Higgs contribution is the same for both the wino and bino. For simplicity, we take the mass of the lightest Higgs boson to be zero in the B functions (we do the same for the wino/bino again) and assume $m_H = m_{H+} = m_A$:

$$\left(\frac{\Delta m_i}{m_i}\right)_{Higgs} = -\frac{\alpha_i}{4\pi}\left[B_1(0,\mu,m_A) + B_1(0,\mu,0) + \frac{\mu}{m_i}\sin 2\beta\left(B_0(0,\mu,m_A) - B_0(0,\mu,0)\right)\right] \tag{3.16}$$

Evaluating the B functions, we find

$$B_1(0,\mu,0) = \frac{1}{2}\left(\ln\frac{Q^2}{m_i^2} - \ln\frac{\mu^2}{m_i^2} + \frac{3}{2}\right) \tag{3.17}$$

$$B_1(0,\mu,m_A) = \frac{1}{2}\left[\ln\frac{Q^2}{m_i^2} - \ln\frac{m_A^2}{m_i^2} + \frac{1}{2} + h(m_A^2/\mu^2)\right] \tag{3.18}$$

where $h(x) = \frac{1}{1-x}\left(1 + \frac{\ln x}{1-x}\right)$. Note that $|h(x)|$ is a monotonically decreasing function of x. To estimate its maximum value, recall that $m_A^2 = 2B\mu/\sin 2\beta$, so taking $B\mu \sim \mu^2$ and setting $\tan\beta = 1$ implies $x = 2$, where $|h(x)| \sim 0.3$. This is already small, and its effect on the spectrum is negligible. To keep consistent with the existing literature, we drop this term. Also,

$$B_0(0,\mu,0) = \ln\frac{Q^2}{m_i^2} - \ln\frac{\mu^2}{m_i^2} + 1 \tag{3.19}$$

$$B_0(0,\mu,m_A) = \ln\frac{Q^2}{m_i^2} - \ln\frac{m_A^2}{m_i^2} + \frac{\mu^2}{\mu^2 - m_A^2}\ln\frac{m_A^2}{\mu^2} + 1. \tag{3.20}$$

3.3.3 NLO Formulae

Adding in the two-loop beta function contribution to m_3, we arrive at the full NLO result

$$\begin{aligned} M_3 &= m_3(Q)\left(1+\frac{3\alpha_3}{4\pi}\left[\ln\frac{Q^2}{m_3^2}+f(r)-\frac{14}{9}\right]+\frac{3\alpha_t}{2\pi}\right) && (3.21)\\ f(r) &= 1+2r+2(r-1)^2\ln|1-r|+2r(2-r)\ln r. && (3.22)\end{aligned}$$

The NLO bino mass is

$$\begin{aligned} M_1 &= m_1(Q)\left(1+\frac{\alpha_1}{8\pi}\left[-22\ln\frac{Q^2}{m_1^2}+11\ln\frac{m_{\tilde{q}}^2}{m_1^2}+9\ln\frac{m_{\tilde{l}}^2}{m_1^2}+\ln\frac{\mu^2}{m_1^2}+\ln\frac{m_A^2}{m_1^2}\right.\right.\\ &\left.\left.+\frac{2\mu}{m_1}\sin 2\beta\frac{m_A^2}{\mu^2-m_A^2}\ln\frac{\mu^2}{m_A^2}-12\right]+\frac{22\alpha_3}{33\pi}-\frac{13\alpha_t}{66\pi}\right) && (3.23)\end{aligned}$$

and the NLO wino mass is

$$\begin{aligned} M_2 &= m_2(Q)\left(1+\frac{\alpha_2}{8\pi}\left[-2\ln\frac{Q^2}{m_2^2}+9\ln\frac{m_{\tilde{q}}^2}{m_2^2}+3\ln\frac{m_{\tilde{l}}^2}{m_2^2}+\ln\frac{\mu^2}{m_2^2}+\ln\frac{m_A^2}{m_2^2}\right.\right.\\ &\left.+\frac{2\mu}{m_2}\sin 2\beta\frac{m_A^2}{\mu^2-m_A^2}\ln\frac{\mu^2}{m_A^2}+1.2+4.32\ln\left(\frac{m_2}{m_W}-0.8\right)\right]\\ &\left.+\frac{6\alpha_3}{\pi}-\frac{3\alpha_t}{2\pi}\right). && (3.24)\end{aligned}$$

Notice that while our expression for the NLO gluino mass agrees with [45], there is a difference in the wino and bino masses, most notably in the coefficients of the $\ln Q^2$ terms. It seems that the authors of [45] did not include a part of the Higgsino/Higgs loop contribution in case of the bino, and omitted gauge boson loop contributions to the wino. Although numerically small, these are required for theoretical consistency, i.e. so that $\frac{dM}{d\ln Q} = 0$ at one-loop order, after plugging in for the one-loop running of the gauge couplings. For the sake of completeness, we have also included a two loop α_2 contribution to the wino mass, which although it is small, should not in principle be dropped because it is of roughly the same size as the other terms.

Finally, for the case of a wino LSP, we consider the splitting between the lightest chargino, i.e. $\widetilde{W^{\pm}}$, and the LSP, $\widetilde{W^0}$, which is important for understanding the phenomenology of the gluino cascade decay. Since the tree-level splitting due to mixing in the neutralino and chargino mass matrices is small for moderate to large μ, the dominant contribution turns out to be due to gauge boson loops:

$$\Delta m_{\widetilde{\chi}} = m_{\widetilde{\chi}^+} - m_{\widetilde{\chi}^0} = \frac{\alpha_2}{2\pi}\left[-2B_0(m_2, m_2, m_W) + B_1(m_2, m_2, m_W)\right]. \tag{3.25}$$

This is more conveniently expressed as

$$\begin{aligned} \Delta m_{\widetilde{\chi}} &= \frac{\alpha_2 m_2}{4\pi}\left[f(r_W) - \cos^2\theta_W f(r_Z) - \sin^2\theta_W f(0)\right], & (3.26)\\ f(r_i) &= \int_0^1 dx\,(2+2x)\ln[x^2 + (1-x)r_i^2], \qquad r_i = \frac{m_i}{m_2}. & (3.27)\end{aligned}$$

The splitting is roughly independent of m_2: For $m_2 = 260$ GeV, $\Delta m_{\widetilde{\chi}} = 167$ MeV, while for $m_2 = 2.6$ TeV, $\Delta m_{\widetilde{\chi}} = 172$ MeV.

3.4 Messengers and Sleptons

To solve the tachyonic slepton problem, we introduce vector-like messenger fields in complete representations of $\mathbf{5}+\bar{\mathbf{5}}$, which get masses from the following tree-level Kähler potential terms:

$$\lambda \int d^4\theta \frac{\phi^\dagger}{\phi} \mathbf{5}\bar{\mathbf{5}} + \kappa \int d^4\theta \frac{X^\dagger X(\mathbf{5}^\dagger \mathbf{5} + \bar{\mathbf{5}}^\dagger \bar{\mathbf{5}})}{M_*^2} \tag{3.28}$$

where ϕ is the conformal compensator, X is the hidden sector field that breaks SUSY, and M_* is a UV cutoff that is naturally the Planck scale in our model. In general we consider N such sets of messenger fields, where we need $N \leq 4$ to preserve gauge coupling unification. Unification also works with one set of $\mathbf{5}+\bar{\mathbf{5}}$ and one set of $\mathbf{10}+\overline{\mathbf{10}}$ (we retain gauge coupling perturbativity in this case because the messengers have masses above 5 TeV [47]). In fact, this is what we need for our complete model with the μ problem solution, to be described below. Since a $\mathbf{10}$ has a Dynkin index of 3/2, this gives the same contribution to soft masses as a model with $N=4$, as is clear from the formulae below.

Since it would be overly contrived to now introduce another scale in addition to $m_{3/2}$, we give the messengers masses of this order in a simple way through the EAM approach outlined in [40], which uses the Giudice-Masiero term [41]. When supersymmetry is broken, the compensator acquires an F term VEV so that $\phi = 1 + m_{3/2}\theta^2$. The messengers also get soft masses through the second term in Eq.

(3.28). Since $m_{3/2} \sim F_X/M_{Pl}$, the soft masses are also set by $m_{3/2}$. We parametrize the soft mass in terms of the supersymmetric mass M as

$$m_{soft}^2 = -c^2 M^2, \tag{3.29}$$

where the reason for the minus sign will become clear in a moment. To prevent the breaking of $SU(3)$, we require $c \leq \sqrt{1 - 1/\lambda}$, with $\lambda \geq 1$. For simplicity, we assume the same GM coupling for all generations of messengers and equal soft masses for each messenger pair. Generalizing these assumptions does not significantly change the picture. We therefore have a hybrid theory of gauge and anomaly mediation, with messenger scale $M = \lambda m_{3/2}$, and $F = -\lambda m_{3/2}^2$. The scalar-messenger mass matrix is

$$M^2 \begin{pmatrix} \mathbf{5}^\dagger & \bar{\mathbf{5}} \end{pmatrix} \begin{pmatrix} 1 - c^2 & -F/M^2 \\ -F/M^2 & 1 - c^2 \end{pmatrix} \begin{pmatrix} \mathbf{5} \\ \bar{\mathbf{5}}^\dagger \end{pmatrix} \tag{3.30}$$

In principle, one should also include contributions to F coming from contact terms with X. However, for the sake of simplicity, we assume that the only contribution is from the GM term, which captures the qualitative features of the general corrections. We examine the effect of the messengers on the gaugino spectrum in the next section.

Here we focus on the soft masses, which give rise to Poppitz-Trivedi D-type gauge mediation. This mechanism was used to solve the tachyonic slepton problem in [36,37, 39]. The idea is simple: since the scalars and messengers share gauge interactions, soft masses for the messengers will induce scalar masses. This contribution was calculated by Poppitz and Trivedi [38] to be

$$\Delta m_i^2 = -\sum_a \frac{g_a^4}{128\pi^4} S_M \, C_{ai} \, Str M_{mess}^2 \log\frac{\Lambda^2}{m_{IR}^2}, \tag{3.31}$$

where S_M is the Dynkin index of the messenger field, C_{ai} is the quadratic Casimir, and Λ, m_{IR} are UV and infrared cutoffs, respectively. The logarithm is large, since we take Λ at the GUT scale and $m_{IR} = M$ (The natural cutoff for our model is the Planck scale, so this is not the entire contribution. There is also a correction coming from physics between the GUT and Planck scales that is not log-enhanced, and includes unknown threshold corrections at the GUT scale). As pointed out in [37], the large logs can be resummed using the one-loop RGE's for the gauge couplings, yielding

$$\Delta m_i^2 = -\sum_a \frac{S_M \, C_{ai} \, Str M_{mess}^2 [g_a^2(\Lambda) - g_a^2(m_{IR})]}{8\pi^2 b_a}, \tag{3.32}$$

with b_a the β function coefficient above the messenger scale. In our case, $Str M_{mess}^2 = 4m_{soft}^2$, so we need m_{soft}^2 to be negative in order to get a positive contribution, hence the minus sign in the definition above. Since $\Delta m_i^2 \sim (\Delta g^2/16\pi^2)(c\lambda)^2 m_{3/2}^2 N$, it is clear that this contribution can easily be as large as that from AMSB for relatively small values of $c\lambda$, i.e. they need not even be $O(1)$, pushing the slepton masses positive at the messenger scale.

3.5 Gaugino Spectrum

We now take a closer look at how the messengers change the gaugino spectrum. Integrating out the messengers takes the soft terms off of the anomaly mediated trajectory. This means that to get the gaugino masses at the weak scale, we need to compute them first at the messenger scale and then run down. We first run up the gauge couplings at two loops to do this. Immediately above the messenger scale the gaugino masses are on the anomaly-mediated trajectory. Using the two-loop beta functions, which include contributions from the messengers, we find

$$\begin{aligned}
m_1(M) &= \frac{\alpha_1}{4\pi}\left(11+\frac{5}{3}N+\frac{\alpha_2}{4\pi}(3N+9)+\frac{\alpha_3}{4\pi}\left(\frac{32}{9}N+\frac{88}{3}\right)-\frac{13\alpha_t}{6\pi}\right)m_{3/2} \\
m_2(M) &= \frac{\alpha_2}{4\pi}\left(1+N+\frac{\alpha_2}{4\pi}(7N+25)+\frac{24\alpha_3}{4\pi}-\frac{3\alpha_t}{2\pi}\right)m_{3/2} \\
m_3(M) &= \frac{\alpha_3}{4\pi}\left(-3+N+\frac{9\alpha_2}{4\pi}+\frac{\alpha_3}{4\pi}\left(\frac{34}{3}N+14\right)-\frac{\alpha_t}{\pi}\right)m_{3/2}. \qquad (3.33)
\end{aligned}$$

We work to two-loop order because there is the possibility of near degeneracy of gaugino masses and we are interested in very small gluino-LSP mass splittings, so corrections at the percent level are important. Since $\alpha_3 \sim 3\alpha_2$ at the weak scale, we see that we can make the gluino and wino masses approximately equal at the messenger scale by choosing $N = 2$. There are also threshold corrections from integrating out the messengers, and the exact expression depends on whether or not the messengers have soft masses. Here we first consider the more general case of messenger soft masses, which are needed in our model. We then discuss the simpler case with no soft masses.

3.5.1 Soft Masses

Adapting the formula in [38], we find

$$\Delta m_i = \frac{\alpha_i}{2\pi} f_t(y_1, y_2) N m_{3/2}, \quad f_t(y_1, y_2) = \frac{y_1 \log y_1 - y_2 \log y_2 - y_1 y_2 \log(y_1/y_2)}{(y_1 - 1)(y_2 - 1)(y_2 - y_1)} \tag{3.34}$$

with $y_1 = M_1^2/M^2$, $y_2 = M_2^2/M^2$, where $M_{1,2}^2$ are the eigenvalues of the scalar messenger mass-squared matrix (we adopt a convention where M_1 is the larger of the two). For the simplest case of universal GM couplings and soft masses that we consider

$$y_1 = 1 - c^2 + 1/\lambda, \qquad y_2 = 1 - c^2 - 1/\lambda. \tag{3.35}$$

For λ not too close to 1 and small soft masses, $f_t \sim 0.5$. As $\lambda \to 1$ (soft masses still small), $f_t \to 0.7$ since there is an enhancement due to contributions from higher order terms in F/M^2 [48]. To examine the role of the soft masses, it is useful to consider the effective number of messengers N_{eff}, a continuous variable defined by

$$N_{eff} = \frac{1 + 2f_t}{2} N. \tag{3.36}$$

Here we consider both positive and negative m_{soft}^2 for λ not too close to 1. In the positive case, as can be seen in Fig. 3.1, N_{eff} decreases with increasing c. As c becomes

Figure 3.1: Effect of c for positive m^2_{soft}

bigger the soft mass dominates over the supersymmetric mass and $N_{eff} \to N/2$. Note that with positive m^2_{soft}, c is not constrained to be smaller than one. N_{eff} increases with c for negative m^2_{soft}, although it is very gradual until $c \sim 0.6$, as shown in Fig. 3.2.

3.5.2 No Soft Masses

The compression of the gaugino spectrum due to messengers is an interesting aspect of extended AMSB theories that, according to our knowledge, has not been investigated previously. We therefore consider it now as an independent module that can be incorporated into other models, i.e. we just focus on the gaugino masses. In this case there is no reason to keep the messenger soft masses, so we simplify the setup by just keeping the GM term for the messengers. In this case, the threshold

Figure 3.2: Effect of c for negative m^2_{soft}

correction from integrating out the messengers takes the simpler form [43]

$$\Delta m_i(M) = -\frac{1}{g_i}(\beta_i' - \beta_i)G(F/M^2)\frac{F}{M}, \tag{3.37}$$

where β' is the beta function above the messenger threshold, β is the beta function below the messenger threshold, and G is the enhancement factor mentioned above. G(x) increases monotonically from 1 to 1.386 as x goes from 0 to 1 [48]. Here we also take different couplings λ_D and λ_T for the triplets and doublets, although the coupling cancels out in F/M, which means that we can only adjust the value of the higher order contribution G separately for the triplets and doublets. Explicitly, the threshold corrections to two-loop order are

$$\Delta m_1 = \frac{\alpha_1}{4\pi}\left(\frac{5}{3}+\frac{3\alpha_2}{4\pi}+\frac{8\alpha_3}{9\pi}\right)\left[G(x_D)+\frac{2}{3}G(x_T)\right]\frac{3}{5}Nm_{3/2} \tag{3.38}$$

$$\Delta m_2 = \frac{\alpha_2}{4\pi}\left(1+\frac{7\alpha_2}{4\pi}\right)G(x_D)Nm_{3/2} \tag{3.39}$$

$$\Delta m_3 = \frac{\alpha_3}{4\pi}\left(1+\frac{17\alpha_3}{6\pi}\right)G(x_T)Nm_{3/2}, \tag{3.40}$$

where $x = -1/\lambda$.

We then run down the gaugino masses to 1 TeV using the one-loop RGEs (we include the next-to-leading order correction for the gluino) and compute the pole masses by adding the corrections appearing in the square brackets in Eqs. (3.21), (3.23), and (3.24). To keep the log term from getting too large, we compute the gluino pole mass at a separate scale, equal to the running mass at 1 TeV. Here we use the $\overline{MS}$ equations so the pole mass equation for the gluino is modified to

$$M_3 = m_3(Q)\left(1+\frac{3\alpha_3}{4\pi}\left[\ln\frac{Q^2}{m_3^2}+f(r)-1\right]\right). \tag{3.41}$$

Finally, to calculate the splitting between the gluino and the LSP, we diagonalize the neutralino mass matrix using the calculated pole masses for the bino and wino.

We now discuss the gaugino spectra for different numbers of messengers. Examples are presented in Table 3.1, with M_3 the gluino mass and ΔM the gluino-LSP splitting. We assume that the scalars and the Higgsinos are somewhat heavier than the gluino; here we choose an arbitrary mass of 1.6 TeV for all, and take tan $\beta = 5$. With heavy squarks, the dominant SUSY production mechanism at the LHC is gluino

N	$m_{3/2}$ (TeV)	λ_D	λ_T	M_1 (GeV)	M_2 (GeV)	M_3 (GeV)	ΔM (GeV)
1	70	2.5	2.5	851	618	523	$\tilde{g}$ lightest
2	40	2.5	1.2	608	580	650	74
3	40	1.5	4.0	721	818	1130	411
4	34	1.5	4.0	712	888	1439	728

Table 3.1: Gaugino spectra

pair production. Each gluino then eventually decays to the bino or wino LSP. In either case, direct decay to the LSP, $\widetilde{g} \to jj\widetilde{\chi}^0$, is possible and is the dominant mode for a bino LSP. For the case of a wino LSP, there is also cascade decay through the charged wino, $\widetilde{g} \to jj\widetilde{\chi}^{\pm}$. For $\Delta m_{\widetilde{\chi}} > m_\pi$, $\widetilde{\chi}^{\pm} \to \pi^{\pm}\widetilde{\chi}^0$ happens 98% of the time. For the range of wino masses that we consider $\Delta m_{\widetilde{\chi}}$ is roughly 170 MeV, so this mode is always open. Although the charged wino can travel a macroscopic distance before decaying [45] (about 1 cm in our case), a displaced vertex analysis is not possible because the pion is too soft. Thus, in terms of observable signatures, we can simply describe the decay to a wino LSP as also being direct.

As pointed out in [49], despite the large production cross-section, these events are difficult to detect when the gluino is nearly degenerate with the LSP since the jets from the decay are very soft. Furthermore, these events may not even have the large E_T^{miss} that is usually a hallmark of R-parity conserving theories, meaning that they will be hidden in QCD background. Even if the gluinos are strongly boosted the LSP

Figure 3.3: Squark contribution to gluino mass

momenta will approximately cancel unless the gluino momenta are unbalanced by the emission of initial or final state radiation.

For $N = 1$, we do not obtain an acceptable spectrum–the gluino is always the lightest. This is because in this case the contributions to the gluino mass from anomaly mediation and the messenger threshold correction have opposite sign. Although the correction due to squarks is sizable, bumping up the gluino substantially would require very heavy squarks. This is because the squark correction increases slowly as the ratio of the squark mass to m_3 is raised, as can be seen in Fig. 3.3.

For $N = 2$, the wino is the LSP unless λ_D gets very close to 1, in which case it is the bino. The spectrum is very compressed, with ΔM no bigger than about 80 GeV. As can be seen in Fig. 3.4, ΔM increases with λ_D because the messenger thresh-

Figure 3.4: Gluino-LSP Splitting for $N = 2$

old correction to the wino gets smaller, decreasing its mass. Conversely, decreasing λ_D raises the wino mass more than the bino mass because of the smallness of α_1, eventually pushing the wino above the bino for $\lambda_D \sim 1$.

The $N = 3$ and $N = 4$ cases are not as compressed but are still squeezed substantially compared to pure AMSB. The bino is the LSP in both cases. We focus on gluino masses below 1.5 TeV, since this is the region that many claim has already been ruled out. Varying λ_D has only a slight effect because ΔM is a few hundred GeV whereas Δm_1 is about 200 GeV, meaning that ΔM can only change by a few tens of GeV.

In summary, for 2, 3, or 4 messenger pairs, relatively light gluinos in the range of 600 to 1400 GeV are still viable. Additionally, compression of the gaugino spectrum is attractive in Split Supersymmetry, where the gluino is perhaps the only sign of new physics and is out of the reach of the LHC with the usual AMSB mass hierarchy and wino/bino thermal dark matter. A detailed discussion about messengers in Split SUSY can be found in Ch. 4.

3.6 Complete Example Spectra

In the previous sections we have discussed separately the solution to the tachyonic slepton problem through D-type gauge mediation and the compression of the gaugino spectrum from the messenger threshold in EAM. We now combine the two to produce complete spectra for different numbers of messengers in the theory. These are listed in Table 3.2. We assume that a solution to the μ problem exists (more on this in the next section) and take $\mu = 1$ TeV (with $B\mu \sim \mu^2$) and three values of $\tan\beta$. For these chosen parameters, we calculate the additional contributions to the Higgs soft masses that are needed for proper EWSB and add them at the messenger scale.

In the $N = 2$ and $N = 3$ cases, we choose squarks in the range of $1.7 - 1.9$ TeV, with a slightly lighter, right-handed stop (there is little mixing), and sleptons that are lighter by about a factor of two. The $N = 4$ case has slightly heavier squarks compared to these. In principle, the squarks could be much lighter in these scenarios, and sneutrinos would become the LSP (a spectrum we do not study, as we are investigating

		$N = 2$	$N = 3$	$N = 4$
inputs:	$m_{3/2}$	40000	40000	30000
	λ	1.3	2.5	3.0
	c	0.25	0.10	0.09
	$\tan\beta$	10.0	13.0	3.0
	μ	986	984	979
sleptons:	$m_{\tilde{e}_L}$	750	762	975
	$m_{\tilde{e}_R}$	657	754	669
	$m_{\tilde{\nu}_L}$	750	762	975
squarks:	$m_{\tilde{u}_L}$	1880	1959	2254
	$m_{\tilde{u}_R}$	1704	1752	2079
	$m_{\tilde{d}_L}$	1880	1959	2254
	$m_{\tilde{d}_R}$	1732	1805	2060
stops:	$m_{\tilde{t}_1}$	1532	1565	1884
	$m_{\tilde{t}_2}$	1816	1889	2177
gauginos:	$m_{\tilde{B}}$	619	714	645
	$m_{\tilde{W}}$	616	795	795
	$m_{\tilde{g}}$	703	1193	1439
Higgs sector:	m_A	5727	5715	3417

Table 3.2: Example spectra for $N = 2, 3, 4$. All masses are in GeV.

scenarios with lighter gauginos). Since the D-type GMSB contribution dominates over its AMSB counterpart – it is about an order of magnitude larger – the mass hierarchy can be explained by ratios of gauge couplings, as in minimal GMSB. However, the near equality of the slepton masses in the $N = 2$ and $N = 3$ cases is a departure from this and arises from both EWSB constraints and our choice of larger $\tan\beta$. Because μ is fairly big, we need $m_{h_d}^2$ to be large to obtain sizable values of $\tan\beta$. In addition, $m_{h_u}^2$ must be negative to satisfy the other EWSB equation involving m_Z. This means that the oft-neglected $\alpha_1^2(m_{h_u}^2 - m_{h_d}^2 + ...)$ piece in the RGE for $m_{\tilde{e}_R}^2$ is significant and drives the mass up considerably when running down. In the $N = 4$ case, we chose a smaller $\tan\beta$ so $m_{h_d}^2$ is not as large and $m_{h_u}^2$ is positive, and this effect is greatly diminished.

The gaugino masses are similar to those in Table 3.1 because the messenger soft masses have only a slight effect on the messenger threshold correction for the small values of c that we need. As noted before, we ignore the $N = 1$ case since that results in a gluino LSP. The mechanism that generates $\mu/B\mu$ will obviously have an impact on the physical Higgs mass, which does not concern us here.

3.7 Discussing the μ Problem

It is well known that one cannot write down a tree-level μ term in minimal AMSB because the resulting $B\mu$ would be a loop factor too large to allow for proper EWSB.

However, D-type gauge mediation is an extra element in our model that contributes to the Higgs soft masses. Taking inspiration from [50], we examine whether we can use this extra freedom to increase the Higgs soft masses enough to make EWSB work despite the large $\mu/B\mu$ hierarchy. We find roughly that to barely satisfy the EWSB stability condition ($\tan\beta = 1$), $c\lambda \sim 3$, yielding squark masses close to 15 TeV, which is clearly not acceptable. Although the two Higgs doublets can be considered a "messenger pair", a GM term is not an option; since μ is large, obtaining a light Higgs would require fine tuning. It is clear that we must add something new to our model.

We can try to solve the μ problem by taking the simplest extension, one that was considered in the early days of gauge mediation [51]. We generate $\mu/B\mu$ by coupling the messengers to the Higgs doublets in the following way:

$$W \supset z_u \overline{\mathbf{10}}\,\mathbf{5} H_u + z_d \mathbf{10}\,\bar{\mathbf{5}} H_d. \tag{3.42}$$

This again leads to the same $\mu/B\mu$ hierarchy, but also gives an extra contribution to the Higgs soft masses, so that we don't have to rely solely on the GMSB contribution. Working to all orders in F/M^2 and including the soft masses, the new yukawa couplings generate the following contributions to the Higgs sector:

$$\mu = \frac{z_u z_d}{2\pi^2} f_t(y_1, y_2) m_{3/2} \tag{3.43}$$

$$B\mu = -\frac{z_u z_d}{4\pi^2} \ln(1-x^2)\lambda^2 m_{3/2}^2 \tag{3.44}$$

$$\Delta m_{h_{u,d}}^2 = -\frac{z_{u,d}^2}{4\pi^2}\left[\left(2-c^2\right)\ln(1-x^2) + \left(1-c^2\right) x \ln\frac{1+x}{1-x}\right]\lambda^2 m_{3/2}^2, \tag{3.45}$$

$$x = \frac{1}{\lambda(1-c^2)}. \tag{3.46}$$

For illustrative purposes only, we also give the formulas for the soft parameters to lowest order in $1/\lambda^2$ and c^2. These should roughly show the correct qualitative behavior since $c < 1$ and $\lambda \geq 1$:

$$B\mu = \frac{z_u z_d}{4\pi^2}(1+2c^2) m_{3/2}^2 \tag{3.47}$$

$$\Delta m_{h_{u,d}}^2 = \frac{z_{u,d}^2}{4\pi^2} c^2 m_{3/2}^2. \tag{3.48}$$

As first noted in [51], it is clear from the above that the Higgs soft masses do not get a contribution at lowest order with messenger soft masses equal to zero. Because of this fact, we must introduce a hierarchy between z_u and z_d so that $\tan\beta$ is not fixed too closely to one, i.e. we need $m_h^2 > B\mu$ for one of the Higgs doublets. So the EWSB will be achieved using the approach outlined in [50], with a large $\mu/B\mu$ hierarchy and a smaller one between $B\mu$ and either $m_{h_u}^2$ or $m_{h_d}^2$. Taking this to be h_u, we find that $\sin 2\beta \sim 2(z_d/z_u)$. Since we do not want Higgsinos that are too light, we must fix the product $z_u z_d$. However, we cannot make z_u larger than 1 without hitting a Landau pole before the GUT scale. These constraints imply that μ is small, about 300 GeV.

The gluino mass with a scalar spectrum similar to that of the $N = 2$ or $N = 3$ cases in Table 2 is then about 1.5 TeV. Although this spectrum is viable (barely), it is only because of the size of the gluino mass and not because of compression, so we do not find this attempt at a μ problem solution to be satisfactory in this minimal form. Note also that we need to tune the value of the other yukawa coupling so that the condition for EWSB in this case, $B\mu^2 > m_{h_u}^2 m_{h_d}^2$, nearly becomes an equality. This has to be done so that we obtain the experimental value of m_Z in the other EWSB equation.

A further issue is the physical Higgs mass. In our model with all messenger soft masses negative, the scalars in each messenger multiplet are lighter than the fermions, so the new Higgs couplings will produce a negative contribution to the Higgs mass squared, whereas to achieve a Higgs mass of 126 GeV we need a large, positive correction for our case of $\sim$ 2 TeV stops [21]. It seems this can be overcome by choosing different soft masses for the **5** and **10** multiplets, with the scalar soft masses of the **10** negative so that the D-type GMSB contribution to the MSSM scalar masses is still positive. Analogously to the top/stop contribution, the messenger contribution to the Higgs mass depends on the logarithm of the ratio of the geometric average of the scalar masses to the fermion mass. So if the scalars from the **5**'s have positive soft masses that are larger in magnitude than those of the scalars from the **10**'s, the average messenger scalar mass can be made larger than the fermion mass, yielding a positive contribution to the Higgs mass.

To obtain a viable solution, we clearly need an additional contribution to μ. This can be done by introducing a singlet and working in the context of the NMSSM. This has been investigated recently in the context of GMSB with new messenger couplings in [52]. In this case the physical Higgs mass results from the large A terms that are produced in this model. Since the singlet is not charged under the SM gauge groups, it does not affect the AMSB or GMSB contributions to the soft masses. In particular, the gaugino spectrum remains unchanged. This model could then eventually be realized in a 5D brane-world setup similar to that of [36, 37], where the visible (MSSM) and hidden sectors are localized on different branes, with the gauge fields and the messengers in the bulk.

3.8 Conclusions

The most recent data from the LHC excludes gluinos with masses less than ~ 1.5 TeV in typical models that have a significant gluino-LSP mass splitting, putting a strain on naturalness. However, gluinos as light as 550 GeV are still allowed for very small mass splittings. We have presented a simple and novel way that such a compressed gaugino spectrum occurs naturally in the context of AMSB.

AMSB models typically have spectra that feature a large spread in the gaugino masses – the gluino is almost 10 times heavier than the wino LSP. Such models can be out of the reach of the LHC for LSP masses of $\mathcal{O}(1 \text{ TeV})$. We find that the presence of messengers in the AMSB framework compresses the spectrum. The resulting spectra have a relatively light gluino with a mass in the range of 600 to 1400 GeV that is no heavier than about twice the LSP mass, with the exact values dependent on the number of messengers N used. For the $N = 1$ case, the gluino is the LSP, while the $N = 2$ case yields a wino or bino LSP depending on the value of the coupling in the Giudice-Masiero term for the messenger doublets, and the mass splitting between the gluino and the LSP is of the order of tens of GeV. The $N = 3, 4$ cases are less compressed and yield a bino LSP.

We have provided expressions for the gaugino pole masses which differ from the expressions present in the literature for the case of the wino and bino. We would like to emphasize that we have confidence in our expressions being correct as the pole masses are independent of the running scale in our case. We have discussed in detail the steps to obtain the pole masses to account for the differences.

Apart from compressing the spectrum, the messengers are crucial in building a complete phenomenological model without singlets, as they help solve the tachyonic slepton problem in AMSB in a way previously suggested in the literature. Contact terms between the messenger fields and hidden sector SUSY breaking result in soft masses for the messengers. Through Poppitz-Trivedi D-type gauge mediation, these soft masses generate contributions to scalar masses which are positive if the soft masses squared are taken to be negative. This contribution is of the same order or greater than the AMSB contribution and thus solves the tachyonic slepton problem.

The only hurdle to a complete model that then remains is the μ problem. We have made an attempt that indicates that solutions can be found in our framework. It involves new yukawa couplings between the Higgs doublets and messengers (i.e. not new fields). However, the minimal version that we considered is deficient, since it produces a light Higgsino LSP and requires considerable fine-tuning. We leave open for future work possible extensions with both positive and negative messenger soft masses and/or the NMSSM as an avenue for resolving the remaining problems.

Chapter 4

Simply Unnatural Supersymmetry

4.1 Introduction

In this chapter, we describe and explore the simplest picture of the the world arising from fine-tuned supersymmetric theories. Our guiding principle is that the model should be "simply un-natural". There is an explicit, un-natural tuning for the weak scale with a clear "environmental" purpose, but in every other way the theoretical structure should be as simple as possible. To this end, we will follow where the theory leads us, without any clever model-building gymnastics. Following what theories of supersymmetry breaking "want to do" leads us to theories with a "minimally split" spectrum where gauginos are near 1 TeV, while scalars, Higgsinos, and the gravitino are parametrically heavier by a loop factor, at a scale m_{sc} between $\sim 10^2 - 10^3$ TeV. This kind of spectrum has long been a ubiquitous feature of simple, concrete models of SUSY breaking. Its modern manifestation was in the context of theories with anomaly mediated SUSY breaking [13], *without* the clever sequestering mechanism of [12].

In [13], the heavy scalars were thought of as something of an embarrassment. This spectrum was later proposed as a serious possibility for supersymmetric theories in [22, 23, 53], put forward as the "simplest model of split SUSY" in [54], and further studied in [55]. We re-initiated a study of this scenario in [56,57]. For obvious reasons, this spectrum has received renewed attention of late [58–63]. The Higgs mass prefers this "minimally split" spectrum, rather than the more radical possibility of scalars up to around $\sim 10^{13}$ GeV [22]. This is perfectly in line with the "simply un-natural" perspective, since theories with much heavier scalars need extra theoretical structure to suppress gaugino masses by much more than a loop factor relative to the gravitino mass.

With this split spectrum, gaugino masses receive comparable contributions from anomaly mediation and the heavy Higgsinos, as well as other possible vector-like matter near the scale m_{sc}. As we will see, this picture has important consequences for flavor physics, as well as a host of novel collider signals that constrain the scale m_{sc} in an interesting way.

In Sec. 4.2, the Mini-Split framework is presented in some detail showing that it is natural to get Higgsino masses of the same order as that of the scalars, and that this is consistent with unification. There we also introduce new vector-like states and study the effects that such messengers have on unification and the gaugino spectrum, and discuss the implications for dark matter. In Sec. 4.3, we point out that a radiative model of flavor arises simply in the Mini-Split setup. A discussion of the experimental signals of our model is found in Sec. 4.4, and we conclude in Sec. 4.5.

4.2 Simplest Tuned Picture of the World

4.2.1 Model and Spectrum

Supersymmetry breaking must give all superpartners in the MSSM masses above their current bounds. Once supersymmetry is broken, there are no symmetries protecting the sfermion masses, and thus scalar masses are expected at some level. On the other hand, (Majorana) gaugino masses require the breaking of an R symmetry, and are thus not guaranteed to arise at the same level. In the case where supersymmetry breaking is communicated via irrelevant operators suppressed by a scale M, sfermion and gaugino masses could arise from the operators of the form

$$\int d^4\theta \frac{X^\dagger X Q^\dagger Q}{M^2} \quad \text{and} \quad \int d^2\theta \frac{Y W^\alpha W_\alpha}{M}, \tag{4.1}$$

where Q and W^α represent visible sector matter and gauge superfields respectively and X and Y are hidden sector chiral superfields which have non-zero VEVs in their auxiliary ('F') components. There are no requirements (other than the absence of a shift symmetry) on the quantum numbers of X, and thus it could be any field from the hidden sector. On the other hand, Y is required to be an exact gauge and global singlet. This stringent requirement makes it clear that gaugino masses will typically not be the same size as scalar masses for generic hidden sectors. In fact, most models of supersymmetry breaking sectors do not contain such singlets [64–71], and this affects both gravity and gauge mediation. While this problem has been 'solved', in the sense that models generating larger gaugino masses have been found [72, 73]–with non-generic superpotentials and/or many discrete symmetries imposed–we take the position that generic models of supersymmetry breaking produce much larger scalar masses than gaugino masses, that is, this is what the models *want* to do. In line with our "simply un-natural" philosophy, we assume the theory-space tuning required to have degenerate sfermions and gauginos is more severe than that required to get the correct electroweak scale.

Another contribution to superpartner masses that theories of broken supersymmetry "want to generate" arises from anomaly mediation. The breaking of R symmetry associated with tuning away the cosmological constant with a constant superpotential gives rise to gaugino masses of order a loop factor beneath the gravitino mass. While there are clever ways to suppress this contribution [22, 74], we consider this contribution generic. Thus, in gravity mediated theories (where M is approximately the Planck scale or a bit below), the gaugino masses will typically end up a loop below the scalar masses.

In Planck- or string-scale mediation of supersymmetry breaking, one possibility for removing the dominant scalar mass operator in Eq. (4.1) is through sequestering, *i.e.*, separating the visible and hidden sectors in an extra dimension or 'conformal throats' [12, 75]. There has been some debate about how generic such sequestering is ([76–79]). We will make the assumption that sequestering is not generic.

We are led to a class of gravity mediation models in which the gaugino masses are roughly a loop factor below the scalars. A possibility we will not explore is single-sector gauge mediation, where again gaugino masses tend to be much more than a loop-factor lighter than scalar masses.

In addition, for "A terms" to be at the same scale as scalar masses, they would have to be generated by operators like

$$\int d^2\theta \frac{Y H_u Q U^c}{M}, \tag{4.2}$$

again requiring a gauge singlet in the hidden sector. Thus our philosophy suggests A terms are small – again, dominated by a one-loop suppressed contribution from anomaly mediation. This will of course have an important impact on the Higgs mass predictions.

Of course the natural version of these models were ruled out when gauginos were not discovered in the 1 GeV range! If these theories are realized in Nature, some kind of "pressure" on the measure pushing towards higher supersymmetry scales is needed, which counteracts the tuning of the cosmological constant and the electroweak scale. We will not attempt to address the notoriously ill-defined question of quantifying these pressures. We will simply assume that whatever the measure is, the likelihood of having a hidden sector that produces degenerate sfermions and gauginos is much smaller than that of a split spectrum with the obvious fine-tuning for electroweak symmetry breaking. We stress that with the spectra we are considering, the fine-tunings at the $\sim 10^{-4} \to 10^{-6}$ level are obviously very severe from the perspective of naturalness, but are dwarfed by the $10^{-60} - 10^{-120}$ levels of fine-tuning for the cosmological constant or the usual 10^{-30} level tuning for the hierarchy problem.

Higgsinos, the μ term, and the Giudice-Masiero mechanism

What about the Higgsinos? The μ term, $W \supset \mu H_u H_d$, breaks both the Peccei-Quinn (PQ) symmetry and potentially an R symmetry, and thus there can be trivial reasons why it is much smaller than the Planck scale. The simplest operator that generates a μ term is the one suggested long ago by Giudice and Masiero [41]:

$$\lambda \int d^4\theta H_u H_d \tag{4.3}$$

where λ is an arbitrary coefficient. In global (flat-space) supersymmetry, this operator represents a total derivative. When including supergravity using the conformal compensator language [80–82], one should multiply this operator by $\phi^\dagger/\phi$ (assuming conformal weights of fields corresponding to their canonical dimensions). The compensator is given by $\phi \simeq 1 + \theta^2 m_{3/2}$, where $m_{3/2}$ is the gravitino mass, as long as the theory has no Planck scale VEVs [83]. Integrating out the Higgsinos and some of the scalars will generate a gauge-mediated-like contribution to the gaugino masses at one loop. The contribution will take gauginos off of the 'anomaly-mediated trajectory' in a special way – a right-magnitude, but wrong-sign contribution to gaugino masses [40, 42]. However, the threshold correction will be affected by squared soft masses for the scalars, and is suppressed when $m_{sc}^2 > m_{3/2}^2$.

In addition, the operator itself appears highly tuned when seen from a different frame. One can remove a chiral operator $\mathcal{O}$ from the Kähler potential $\mathcal{K}$ via the transformation

$$\mathcal{K} \to \mathcal{K} - (\mathcal{O} + \mathcal{O}^\dagger) \quad \text{and} \quad W \to W e^{\mathcal{O}/M_{pl}^2}. \tag{4.4}$$

For $\mathcal{O} = \lambda H_u H_d$, the term in Eq. (4.3)) becomes terms in the superpotential, $(1 + \lambda(H_u H_d/M_{pl}^2) + ...)(W_{hid} + W_0 + W_{vis})$, where W_{hid} contains the fields involved in dynamical supersymmetry breaking and W_0 is the operator that generates a constant superpotential. Thus a pure Giudice-Masiero (GM) term, Eq. (4.3)), is a result of a precise relationship between the coefficients of two operators, $H_u H_d W_{hid}$ and $H_u H_d W_0$ (up to H^2/M_{pl}^2 corrections). This particular combination could result from a direct coupling of the curvature, whereas direct couplings to the constant superpotential and supersymmetry breaking sectors could be suppressed due to sequestering. However, we are assuming that there is no sequestering, and thus we do not have a predictive relationship between the effective μ and $B\mu$ terms. For example, if only the $H_u H_d W_0$ operator existed, the threshold would be purely supersymmetric, and, in the limit of vanishing scalar soft masses, would keep gauginos on the anomaly-mediated trajectory (of the MSSM without Higgsinos).

Having said this, for the sake of simplifying parameter space, we will take the case of pure GM as the 'central value' of the threshold correction in theory space. Regardless of the details of the threshold, it is clear that it is trivial to generate $\mu \sim m_{3/2}$ in multiple ways.

Of course since the μ term also breaks PQ symmetry, it is possible to imagine that the Higgsinos are lighter, near the same scale as the gauginos, as in the earliest models of Split SUSY [22].

One can imagine suppressing these operators, as they explicitly break the PQ symmetry (under which $H_u H_d$ is multiplied by a phase). For a pure GM term, approximate PQ symmetry implies $\lambda \ll 1$ and for $m_{scalar} \sim m_{3/2}$, this leads to a suppression of $\mu = \lambda m_{3/2}$ and $B\mu = \lambda m_{3/2}^2$, and thus $\mu \sim m_{scalar}/\tan\beta$. In the limit where $\lambda \to 0$, or more generally when Planck-suppressed superpotential couplings between $H_u H_d$ and W_0 or the hidden sector are absent, $\tan\beta$ is large yet the Higgs mass requires low scalar masses, thus rendering the spectrum unviable. In principle, the Kähler potential operator $X_{hid}^\dagger X_{hid} H_u H_d$ could generate $B\mu \sim m_{scalar}^2$, in which case μ would be generated by gaugino loops such that $\mu \sim (\alpha/4\pi) m_{gaugino}$. However, we see no symmetry reason for this limit and do not explore this spectrum further (its phenomenology was explored in [59, 63]).

If $\mu^2 \gg m_{sc}^2$, then electroweak symmetry is only broken if the coefficient of the scalar bilinear $H_u H_d$ (the $B\mu$ term) approaches the limit $B\mu \to \mu^2$. This is an interesting case, as $\tan\beta \to 1$ in this limit, a value which is disfavored in the 'natural' MSSM both because the physical Higgs mass is too low, and because the top Yukawa

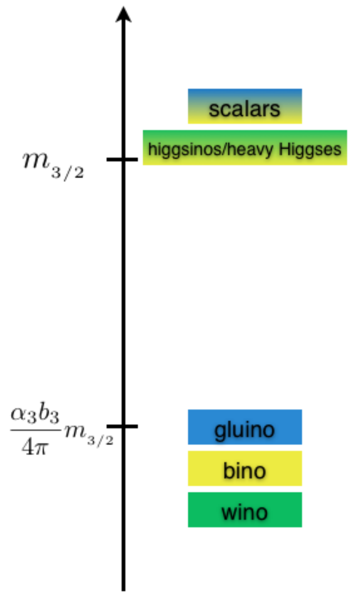

Figure 4.1: A "simply unnatural" spectrum.

coupling gains a Landau pole below the GUT scale. Neither of these present a problem with heavier scalars. However, this limit requires not only the tuning for electroweak symmetry breaking ($\lambda \to 1$), but also the magic of a pure GM mass. There may not be a clear UV reason to naturally favor this point in parameter space, but it is at least interesting that it is allowed phenomenologically and should be considered.

Spectrum and Unification

For our minimal model, we take scalar masses and Higgsinos to be roughly degenerate $m_{sc} \sim \mu \sim m_{3/2}$ (Figure 4.1), and determine the overall scale favored by a Higgs mass of 125 GeV. The result, shown in Figures 4.2 and 4.3, is somewhat different from that found in [86]. The Higgsinos are not present in the effective theory

Figure 4.2: The Higgs mass predicted as a function of the scalar masses and $\tan\beta$. The bands at $\tan\beta = 1$ and 50 represent the theoretical uncertainty in the top mass and α_s. The gaugino spectrum is that predicted by the anomaly mediated contribution with the gravitino mass $m_{3/2} = 1000$ TeV, resulting in an approximate mass for the LSP wino of $\sim 2.7 - 3$ TeV (which is roughly the mass necessary for the wino to have the correct cosmological thermal relic abundance to be all of dark matter [84]). The μ term is fixed to be equal to the scalar mass – this threshold has a small but non-negligible effect on the Higgs mass relative to the conventional split supersymmetry spectrum [22, 23]. The A-terms are small.

Figure 4.3: The allowed parameter space in the $\tan\beta - M_{sc}$ plane for a Higgs mass of 125.7 ± 0.8 GeV, for $\mu = m_{sc}$. The solid blue lines delimit the 2σ uncertainty. The dashed blue lines show the effect of the 1σ uncertainty in the top mass, $m_t = 173.2 \pm 0.9$ GeV [85]. We take the gaugino spectrum predicted by AMSB (including the heavy Higgsino threshold) with the gravitino mass $m_{3/2} = 500$ TeV, resulting in a wino LSP of 2.6 TeV, and a gluino mass of 14.4 TeV. However, the Higgs mass is highly insensitive to the gaugino spectrum, and a gravitino mass of 50 TeV yields essentially the same plot above.

Figure 4.4: The running of the gauge couplings with scalar masses and Higgsinos fixed at 10^3 TeV. Error bands on α_s are at the three-sigma level according to the Particle Data Group [85]. We use $M_{gluino} = 14.4$ TeV, $M_{wino} = 2.6$ TeV, $M_{sc} = \mu = 10^3$ TeV and $\tan\beta = 2.2$ to generate this plot.

beneath m_{sc}, and thus the positive contribution to the running of the Higgs quartic coupling from Higgsino/gaugino loops in the low-energy theory is absent. This means that $m_H \sim 125$ GeV is consistent with heavier scalars than in the split spectrum considered in [86]. In particular, for moderate $\tan\beta$, scalar masses in the $10^3 - 10^4$ TeV range are perfectly allowed, while with light Higgsinos such heavy scalars are only possible if $\tan\beta$ is tuned to be close to 1. Such heavy scalars naturally avoid all flavor problems, giving another impetus to focus on a spectrum with $\mu \sim m_{3/2}$.

Because the Higgsinos have a significant impact on the differential running of gauge couplings, keeping μ heavy noticeably changes the unification prediction. For example, we see in Figure 4.4 that two-loop running predicts a smaller value of the strong coupling constant $\alpha_s(M_z)$ than what is generally found in the 'natural' MSSM. For example, with $\mu = 1000$ TeV, $\alpha_s(M_z) = 0.113$ for gluinos at 1.5 TeV and $\alpha_s(M_z) = 0.111$ for gluinos at 15 TeV (for $\mu = 100$ TeV, $\alpha_s(M_z) = 0.1185$ and 0.117, respectively). Of course this prediction is affected by unknown threshold corrections at the GUT scale, but the values found here are a bit closer to the world average of $\alpha_s(M_z) = 0.1184(7)$ [85], and are very close to more recent determinations using LEP data [87].

Heavy scalars and Higgsinos do moderately better at b-τ unification than the natural MSSM (Figure 4.5), especially at low values of $\tan\beta$. In addition, for small $\tan\beta$, the top Yukawa runs relatively strong at the GUT scale, and one would naturally expect significant threshold corrections.

In pure anomaly mediation, the gaugino masses are widely split, with the gluino roughly a factor of ten heavier than then wino. This is due to the same accident as the near cancellation of the one-loop beta function of SU(2) in the MSSM. With a pure GM term (ignoring soft masses), the Higgsino threshold increases the wino and bino masses such that the gluino/wino ratio is reduced to roughly a factor of six. An interesting limit occurs if the Higgses are mildly sequestered from W_{hid} such that Planck-suppressed couplings to supersymmetry breaking are absent, but the μ-term comes from $H_u H_d W_0$. In such a limit, the threshold correction suppresses

Figure 4.5: The Yukawa coupling ratio (y_b/y_τ) evaluated at the GUT scale as a function of the scalar mass. The gluino masses corresponding to the light and heavy wino masses are 1.5 TeV and 14.4 TeV, respectively. The bands correspond to tan $\beta = 1.2$ and 50 from top to bottom.

the wino mass, and in fact at leading order in $B\mu/\mu^2$ the wino mass vanishes! Of course, without soft masses, electroweak symmetry breaking at a scale much smaller than $m_{3/2}$ would require $B\mu/\mu^2 \to 1$, in which case the wino retains $\sim 40\%$ of its standard MSSM value. Without sequestering, however, soft masses generally reduce the threshold effect, and the operator $H_u H_d W_{hid}$ adds to the magnitude of the wino mass and thus reduces the large splitting.

4.2.2 New Vector-Like States

As with the μ term, $m_{3/2}$ is a natural mass scale for vector-like states. Additional vector-like states, with big SUSY breaking, can further significantly modify the anomaly-mediated spectrum of gauginos. To preserve gauge coupling unification, we assume that these states are in complete multiplets of $SU(5)$. We have seen in Ch. 3 that in the simple limit that their masses come from a pure GM mass term, they invariably produce a squeezed spectrum among the MSSM gauginos. As defined in Sec. 3.5, the effective number of messengers measures the size of the threshold correction compared to that of one canonical $\mathbf{5}+\bar{\mathbf{5}}$ pair (in standard $SU(5)$ language) with a pure GM mass and no additional scalar soft masses.

A heavy vector-like state whose mass comes only from a superpotential (*i.e.*, supersymmetric) operator would, at leading order in F/M, decouple in such a way as to leave the anomaly-mediated relationships between beta-functions and gaugino masses intact. In the case of a pure GM mass term, the sign of the effective F term for the scalar components of the new states is opposite to that of the superpotential case, and therefore the threshold corrections to gauginos also have sign opposite to the one required to keep the spectrum "anomaly mediated." For one to four sets of vector-like states, this tends to suppress the splitting between the gluino and the wino (or lightest gaugino). For example, with one vector-like state, the one-loop beta-function coefficients above the threshold for $SU(3)$ and $SU(2)$ are $b_3 = (b_3)_{\rm MSSM} + 1 = -2$ and $b_2 = (b_2)_{\rm MSSM} + 1 = 2$, respectively. Below the threshold, the coefficients become

Figure 4.6: The gaugino spectrum as a function of N_{eff} (defined in Sec. 3.5) at two-loop order plus threshold corrections. The other parameters in the model are $m_{3/2} = 70$ TeV, $\tan\beta = 2.2$, and the coefficient for the GM term $\lambda = 1.1$. M_{sc} is again the soft mass for all MSSM scalar superpartners and we set $\mu = M_{sc}$.

$(b_3)_{\text{MSSM}} = -3$ and $(b_2)_{\text{MSSM}} = 1$, while the gaugino masses (at leading order) are proportional to $(b_3)_{\text{MSSM}}+2 = -1$ and $(b_2)_{\text{MSSM}}+2 = 3$. Accounting for the hierarchy in gauge couplings, this renders gluinos and winos roughly degenerate. Generally, the gaugino mass coefficients for N messengers will be $(b_i)_{\text{MSSM}} + 2N$, where i runs over the gauge groups.

More generally, soft masses for the scalar components of the vector-like states will suppress the threshold correction. In the limit of soft masses much larger than the GM mass, the threshold correction goes to zero, and the resulting spectrum becomes

Figure 4.7: The running gauge couplings in the case of $N = 1$ vector-like state (dashed), and $N = 4$ (solid). The scalar masses and Higgsinos are fixed at 10^3 TeV. Error bands on α_s are at the three-sigma level according to the Particle Data Group [85]. We use $M_{gluino} = 14.4$ TeV, $M_{wino} = 2.6$ TeV, $M_{sc} = \mu = 10^3$ TeV and $\tan\beta = 2.2$ to generate this plot.

proportional to $(b_i)_{\text{MSSM}} + N$ – only half the effect. Thus, a more general parameterization of this threshold contribution is $(b_i)_{\text{MSSM}} + 2N_{eff}$, where N_{eff} (defined in Sec. 3.5) is N in the GM limit, and $N/2$ in the limit of large scalar soft masses. In Figure 4.6, we plot the gaugino spectrum as a function of N_{eff}. We see that the ratio of gluino mass to lightest gaugino is always smaller than the pure MSSM case. A squeezed spectrum is of course more hopeful for the discovery of the gluino at a collider for fixed LSP mass.

The vector-like states have a slightly negative effect on unification, as shown in Figure 4.7. For example, if we take the masses of the fermionic components of the messengers to be 10^6 GeV, then for $N = 1$, the central value predicted for $\alpha_s(M_Z)$ is 0.111, while for $N = 4$ it is 0.109. In addition, $b - \tau$ unification is significantly worse. However, if these new states are associated with a model of flavor, Yukawa coupling unification would depend on the full theory.

4.2.3 Dark Matter

One of the compelling motivations for new particles at the weak scale is the idea of WIMP dark matter. In models of the sort we are considering, where R-parity makes the LSP stable, we expect some thermal relic abundance regardless of whether the LSP comprises the majority of the dark matter. And since this is the lightest new particle in the spectrum, it is important to understand what mass it can have.

To begin with, we can consider the conventional anomaly-mediated spectrum, with a wino dark matter candidate [88]. In this case, to achieve the appropriate relic abundance, we require a mass of $\sim 2.7 - 3$ TeV [84]. With conventional anomaly mediation for the gaugino masses, this would make the gluinos inaccessible at the LHC. However, as we have already discussed, with the contributions of the Higgsinos and potentially new vector-like states, the spectrum is naturally squeezed. If it is quite squeezed, it is conceivable that the gluinos will be just at the edge of discovery, even with a thermal relic wino dark matter candidate. Since the direct detection cross section of a pure wino is extremely small [89, 90], below $O(10^{-47}\text{cm}^2)$, discovery via direct detection will be extremely difficult.

However, a number of other options are also possible. With a wino LSP, it may simply be that the dark matter is dominantly composed of something else (e.g., axions). In such a case, the LSP can be quite light (from the perspective of cosmological constraints), and almost any spectrum is open to us, including relatively light ($\sim$ TeV) gluinos. Such a wino could be the dark matter if produced non-thermally (e.g., [45,88,91], or more recently [92]). Indeed, in the context of minimal Split SUSY models, it is reasonable to expect late-decaying moduli to dilute any thermal LSP abundance, with the dark matter being re-populated by modulus decays. This still favors dark matter lighter than the TeV scale to get the correct relic abundance, and can also pleasingly dilute the troublesome axion abundance down to acceptable levels, for axion decay constants almost as high as the GUT scale [55]. If the bino is the LSP, we must rely upon late-time entropy production to dilute away an otherwise highly overabundant relic.

In each case, there remains the prospect for interesting collider signals. For a thermal relic, we must count on a squeezed spectrum, while non-thermal (or non-WIMP dark matter) cases are generally easier to find. Regardless, the appearance of signals at the LHC will possibly point to a non-standard thermal history.

4.3 New Flavor Physics and Radiative Fermion Masses

In our picture, the supersymmetric flavor problem must be solved in a trivial way, and not with ingenious model-building and gymnastics. Without any special structure to the scalar mass matrices, in particular with no mechanism enforcing scalar mass degeneracy, $K - \bar{K}$ mixing and ϵ_K demand that the first two generations of squarks be as heavy as $\sim$ 1000's of TeV. What about the third generation squarks? They can plausibly be comparable to the first two generations, or at most an order of magnitude lighter.

To give a simple example for theories of flavor leading to the second possibility, consider models where the Yukawa hierarchy is explained by the Frogatt-Nielsen mechanism, with the light generations having different charges under anomalous $U(1)$ symmetries [93,94]. The anomalous $U(1)$'s are Higgsed by the Green-Schwartz mechanism, and the gauge bosons are lighter than the UV cutoff (string scale), parametrically by a factor of $\sqrt{\alpha}$. Tree-level exchange of this $U(1)$ gauge boson can give SUSY breaking that dominates over Planck-suppressed soft masses. This gives large, different masses to the first two generations, since they are charged under the $U(1)$, but not to the third generation. With an $\mathcal{O}(1)$ splitting between the first two generation scalars, these soft masses must be in the range of at least 1000's of TeV. Planck-suppressed operators will put the third generation scalars in the range of 100's of

TeV. Note that we can't imagine the third generation much lighter than a factor of ~ 10 compared to the first two generations. In RG running from high scales, two-loop corrections to the third generation soft masses from the first two generations give large negative soft masses that would lead to color breaking [95].

Thus, if we want to trivially solve the flavor problem, the first two generations of scalars should be in the range of 1000's of TeV, while the stops can be at most an order of magnitude lighter, in the range of hundreds of TeV. Note that the Higgs mass constrains the geometric mean of the left and right stop masses, given by $m_{\tilde{t}} = \sqrt{m_{\tilde{q}_3} m_{\tilde{u}_3}}$, and as we have seen, with $\tan\beta \sim \mathcal{O}(1)$, we can have $m_{\tilde{t}} \sim 10^2 - 10^3$ TeV, so this is perfectly consistent with solving the flavor problem. It is also interesting that the scalars can't be much *heavier* without making the Higgs mass too big. Thus, the absence of flavor violation pushes the scalar masses up, but getting the right Higgs mass doesn't allow these masses to get too large, saturating right around a scale of $\sim$ 1000 TeV. It is of course notable that this is just what is expected from the simplest picture of SUSY breaking we have been discussing.

If we have scalars in just this range, with no special effort to suppress flavor violation in the soft terms, we might be sensitive to new flavor violation in future experiments.

There is another interesting observation about flavor, which provides an additional motivation for a "split" spectrum with gauginos lighter than scalars. Let us suppose that there is indeed large flavor violation in the soft masses. Huge FCNCs are an

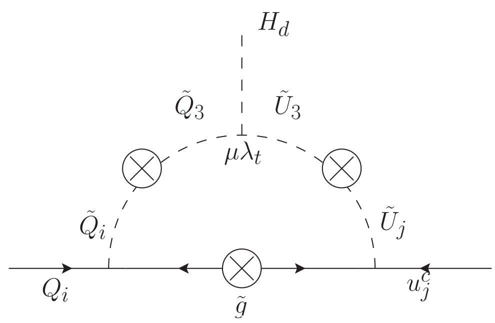

Figure 4.8: Diagram that generates up-type quark Yukawa couplings from the top Yukawa in the case of large mass mixing between flavors, indicated by the crosses on the scalar lines.

obvious worry, but (ignoring detailed issues of the Higgs mass for the moment) one would imagine that these could be decoupled by making the scalars arbitrarily heavy. This is of course correct. But in a theory with no splitting between scalars, Higgsinos, and gauginos, there is a far greater difficulty with flavor that cannot be decoupled by pushing up the scale of SUSY breaking: The large flavor violation, in tandem with a large top Yukawa coupling and the breaking of R-symmetry by the μ term and gaugino masses, radiatively feeds unacceptably large contributions to the up-type quark Yukawa couplings at one-loop. The diagram in Fig. 4.8 yields

$$\delta\lambda_u^{ij} \sim \frac{\alpha_s}{4\pi}\frac{m^2_{\tilde{Q}_{i3}}}{m^2_{sc}}\frac{m^2_{\tilde{U}_{j3}}}{m^2_{sc}}\frac{\lambda_t}{\tan\beta}\frac{\mu\, m_{\tilde{g}}}{m^2_{sc}} \tag{4.5}$$

With $m^2_{ij}/m^2_{sc} \sim O(1)$, $\tan\beta \sim O(1)$, and $\mu \sim m_{\tilde{g}} \sim m_{sc}$, this gives a correction to all up-type Yukawa couplings of order $\sim 10^{-2}$, vastly larger than observed.

It is interesting that for our minimally split spectrum, with $\mu \sim m_{sc}$ and $m_{\tilde{g}} \sim 10^{-2}\, m_{sc}$, this correction is roughly of the order of the observed up quark Yukawa coupling. The up quark mass can plausibly arise from this "SUSY slop." Note that the analogous "slop" cannot be significant for the down and electron Yukawa couplings since the corrections are $\propto \lambda_{b,\tau}$ tan β, and for the moderate tan β we are forced to have, the corrections are about 10^{-2} of the observed values.

More generally, supersymmetric theories with a split spectrum allow us to re-open the idea of a radiatively generated hierarchy for Yukawa couplings. The central challenge to building such theories of flavor is the following: the chiral symmetries protecting the generation of Yukawa couplings must obviously be broken, but then what forces the Yukawas to only be generated at higher loop orders [96, 97]? Supersymmetry offers the perfect solution to this problem, since the chiral symmetries can be broken in the Kähler potential, while holomorphy can prevent this breaking from being transmitted to Yukawa couplings in the superpotential. It is then only after SUSY breaking that the chiral symmetry breaking can be transmitted radiatively to generate Yukawa couplings [98]. Unfortunately, it is extremely difficult to realize this idea in a simple way, with a natural supersymmetric spectrum; the degree of flavor violation needed in the soft terms is large, and would lead to huge flavor-changing neutral currents. But in our new picture this is no longer the case: Yukawa couplings are dimensionless and can be generated at any scale, while the FCNCs decouple as the scalars are made heavy.

As we have seen, with the minimal MSSM particle content, only the top Yukawa coupling is large enough to seed the other Yukawa couplings, and thus it is only possible to generate the up quark Yukawa coupling radiatively. Additional vector-like matter near m_{sc} allows the possibility of new large Yukawa couplings and thus more radiatively generated Yukawas. For instance, with a single additional $\mathbf{5} + \bar{\mathbf{5}}$ messenger pair, $(D^c, L) + (D, L^c)$, we can have large mixing Yukawas of the form

$$\lambda_d^i q_i H_d D^c, \lambda_e^i e_i^c H_d L \tag{4.6}$$

as well as large scalar soft mass mixing

$$m_{d_i}^2 \tilde{D^c}^* \tilde{d^c}^i, m_{l_i}^2 \tilde{L}^* \tilde{l}^i \tag{4.7}$$

between the (D^c, L) and the ordinary d_i^c, l_i. Then, the analogous diagram to Figure 4.8 contributes to down-type quark and charged lepton Yukawa couplings, with $\lambda_t \to \lambda_{d,e}^i$. For $\lambda_{d,e} \sim O(1)$, we can have a radiative origin for the down quark and electron Yukawa couplings.

Given that $m_{\tilde{g}} \sim \frac{\alpha}{4\pi} m_{sc}$, these radiative corrections are parametrically two-loop effects. With additional full vector-like generations, together with heavy-heavy and heavy-light Yukawa couplings to the Higgses, we can also get one-loop corrections through diagrams of the form shown in Figure 4.9. This picture thus easily provides some basic ingredients for the construction of realistic theories where all the fermion masses arise radiatively off the top Yukawa coupling together with other $O(1)$ Yukawa couplings to heavy vector-like states. In Ch. **??**, we build a radiative model of flavor

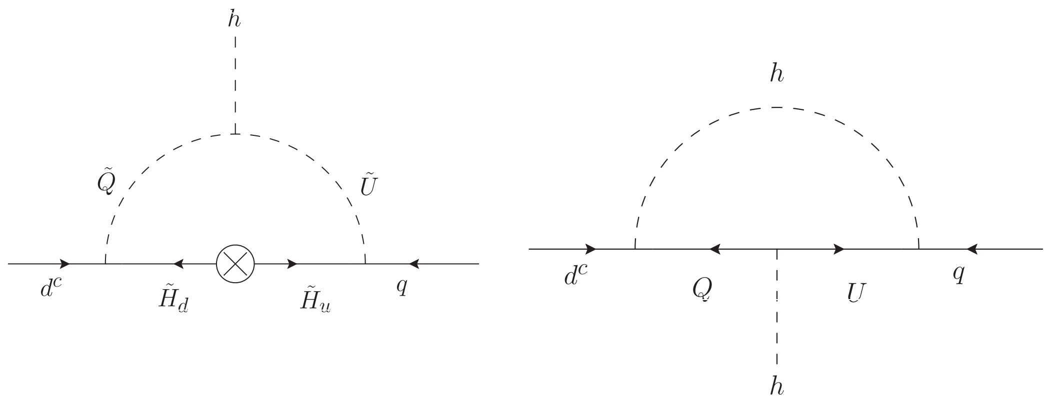

Figure 4.9: Radiatively generated down-type quark Yukawa couplings seeded by heavy messenger-Higgs Yukawa couplings.

along similar lines, with a democratic treatment of all flavors.

4.4 Tests of Un-naturalness

The theoretical developments leading to the development of the Standard Model were greatly aided by concrete experimental evidence for the presence of new physics at short distances not far removed from the scales experimentally accessible at the time. This was most obvious for the weak interactions, which were encoded as dimension six operators suppressed by the Fermi scale. The presence of these operators, together with their $V - A$ structure, were strong clues pointing to the correct electroweak theory. The theoretical triumph of the Standard Model has rewarded us with a renormalizable theory, with no direct evidence at all of higher dimension operators suppressed by nearby scales. Instead of having concrete clues to the structure of new physics through the observation of non-zero coefficients for higher dimension operators–say through a large correction to the S-parameter, a large rate for $\mu \to e\gamma$ or sizable electron EDMs–the main guideline to extending the Standard Model for the past thirty years has been to explain "zero": the *absence* of large quadratically divergent corrections to the Higgs mass, while seeing *no* observable effects in higher dimension operators.

The discovery of a natural supersymmetric theory–as spectacular as it would be–would eventually leave us in a similar position: we would have another renormalizable theory, with no obvious indications for new physics needed until ultra-high energy scales. Amusingly enough, however, the situation is completely different in the un-natural theories we have been discussing in this chapter. The theory has two scales of new physics, $m_{1/2}$ for the gaugino masses and m_{sc} for the scalar masses. As we will see, if we can produce the gauginos, their decays can provide us with unique opportunities to measure or constrain higher dimension operators suppressed by the scale m_{sc}.

Before turning to this discussion, let us first ask an even more basic question: what experimental signals can immediately *falsify* these simply un-natural theories? The existence of any new scalar state beyond the Higgs would immediately exclude simply un-natural models, since it would require an additional tuning, a mechanism to stabilize its mass, or an elaborate family/Higgs symmetry structure. The only important caveat to this is that pseudo-Nambu Goldstone bosons (pNGBs) could still be consistently present. However, if these are light without tuning, they can't be charged under the SM gauge groups, and they can only have higher dimension couplings to SM fields suppressed by their decay constant. We would not expect to produce such states at colliders.

Thus, a second light Higgs, mixing and significantly altering the properties of the Higgs (e.g., [99–101]), would exclude the theories described in this chapter. This makes precision measurements of the Higgs, such as deviations from SM branching ratios, especially important. Large modifications of the Higgs couplings to the W, Z require the existence of a new scalar. Higgs couplings to fermions can be modified by e.g., mixing with new vector-like matter, as can the Higgs width to $\gamma\gamma$. Both ATLAS and CMS reported in their Higgs discovery announcements a ~ 1.7 enhancement of the rate for $h \to \gamma\gamma$ relative to the SM [102–106]. A further CMS analysis based on the full dataset from the first run of the LHC saw this drop to a level that is consistent with the SM to within 1σ [107]; the ATLAS number has though stayed more-or-less constant and disagrees with the SM by about 2σ [108]. Both collaborations however report consistency with the SM in the ZZ and WW decay channels [109–111]. With only extra fermions, it is a challenge in general to theoretically achieve rates for $h \to \gamma\gamma$ above a factor of 1.5 of the SM value, and even then with some tension relative to precision electroweak observables [112, 113]. For the case of Split SUSY, however, getting the enhancement in the $\gamma\gamma$ channel while leaving WW and ZZ unchanged requires something very specific: new electrically-charged vector-like fermions that are light (about 150 GeV or less) and have electroweak interactions but that do not carry color [114]. Thus, if the hint, albeit quite mild at present, of a deviation in $h \to \gamma\gamma$ persists without an associated discovery of charged particles lighter than ~ 150 GeV, simply un-natural theories will have been conclusively excluded.

In fact, within the framework we are discussing, with only the MSSM field content present, the leading interactions of the Higgs bosons to new, electroweak charged states are suppressed by $1/\mu$, and thus the corrections to Higgs properties are far too small to be seen. Thus, at least this minimal framework could be excluded by any convincing deviation of Higgs properties from SM expectations.

4.4.1 Gaugino Decays and the Next Scale

In a simply unnatural theory, the only new particles that we expect to see at the LHC are the gauginos. The impact of this on the decays of gauginos is profound, for a simple reason: *for $SU(3) \times SU(2) \times U(1)$ fermions that do not carry B or L, there are no renormalizable operators under which they can decay into each other and SM particles.* This simple point was emphasized by [22]. As a consequence, the gaugino decays are necessarily suppressed by a new higher scale.

The scale in question varies depending on the particular process. Gluinos decay

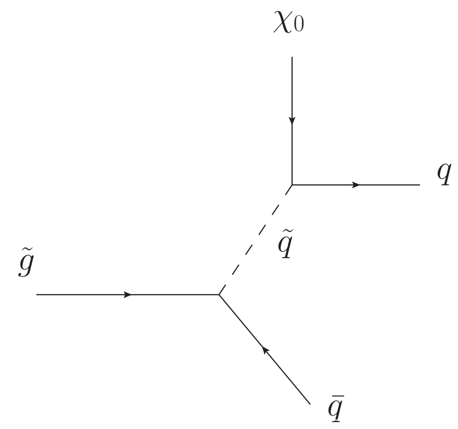

Figure 4.10: Diagram involving a heavy, off-shell squark that yields the dimension-six operator of Eq. (4.8) contributing to gluino decay to the LSP χ^0.

through the diagram in Figure 4.10, which yields the dimension-six operator

$$\frac{g_3 \bar{q}\tilde{g}\bar{\chi}q}{m_{\tilde{q}}^2}. \tag{4.8}$$

The lifetime for such a decay can be quite long, with

$$c\tau \approx 10^{-5}\text{m}\left(\frac{m_{\tilde{q}}}{\text{PeV}}\right)^4\left(\frac{\text{TeV}}{m_{\tilde{g}}}\right)^5 . \tag{4.9}$$

This leads to an interesting immediate observation: the fact that gluinos decay *at all* inside the detector will imply a scalar mass scale within a few orders of magnitude of the gluino mass scale. Moreover, if the gluino decays promptly, without any displacement, we will already know that the scalar mass scale is at an energy scale $\lesssim 100$ TeV, which is at least conceivably accessible to future accelerators.

While this signal places an upper bound on the next mass scale, there are signals that can simultaneously place a quick lower bound. In particular, it is possible to imagine that large flavor violation in the scalar sector could produce clear flavor violation in the gluino decays (e.g., $\tilde{g} \to \bar{t}c$). If so, closing the loop generates sizable flavor-violating four-Fermi operators $\alpha_s^2 q^4/m_{\rm sc}^2$. Even for CP-conserving processes, constraints push this scale to [115] $\sim 10^3\,\text{TeV}$ ($\sim 10^4\,\text{TeV}$ if CP is violated). A combination of a lack of displaced vertices and large flavor violation in gluino decays could quite narrowly place the next scale of physics, without ever having observed a single particle close to the heavy scale.

The quark line above can be closed to yield a chromomagnetic dipole operator as well

$$\frac{g_3^3}{16\pi^2}\frac{m_{\tilde{g}}}{m_{\tilde{q}}^2}\log(m_{\tilde{q}}/m_{\tilde{g}})\tilde{g}^i_j\sigma^{\mu\nu}\tilde{b}G^j_{i\mu\nu}. \tag{4.10}$$

Such an operator will produce dijet + MET (missing transverse energy) signals, but because their rate is suppressed by a loop factor, they should be lost in the overall four jet + MET signals of the off-shell squark decay.

In contrast to gluinos, bino decay proceeds through a dimension-five operator that arises from integrating out the Higgsino, shown on the left in Figure 4.11, namely

$$\frac{g_2 g_1}{\mu}h_i^*\tilde{W}^i_j h^j\tilde{B}\ . \tag{4.11}$$

The suppression by only one power of the heavy scale suggests that these decays will be prompt.

Note that this operator generates a $\tilde{W}-\tilde{B}$ mixing term, which in general will correct the mass of the neutral wino relative to the charged wino by an amount

$\sim m_W^4/(\mu^2 m_{\tilde{W}}) \sim 10^{-7}\text{GeV}(\text{PeV}/\mu)^2(\text{TeV}/m_{\tilde{W}})$. This correction is negligible, even compared to the conventional loop-suppressed mass splitting $m_{\tilde{W}^\pm} - m_{\tilde{W}^0} \approx 150$ MeV [116]. The leading dimension-five operator $\tilde{W}^i_j h^*_i h^l \tilde{W}^j_l$, while correcting the wino mass, does not yield any mass splittings between the usual components. Thus, the mass splitting between charged and neutral winos is a clear test of heavy Higgsinos.

The operator of Eq. (4.11) leads to the decay $\tilde{W}^0 \to h\tilde{B}$. Note, however, that because there is no light *charged* Higgs, there can be no decay $\tilde{W}^\pm \to h^\pm \tilde{B}$ through this operator. Rather, the equivalent decay will arise from the resulting $\tilde{W} - \tilde{B}$ mixing, giving $\tilde{W}^\pm \to W\tilde{B}$. While the $\tilde{W}^\pm \tilde{W}^0$ production cross section is generally smaller than the Higgs production cross section, it is not far off from the associated production cross section (a few hundred fb at $m_{\tilde{W}} \sim 200\,\text{GeV}$ [117]). Consequently, in these models, there are new avenues for Wh production that might be searched for.

Note that this dimension-five operator does *not* give $\tilde{W}^0 \to \tilde{B}Z$. This decay can arise from a dimension-six operator, integrating out the Higgsinos at tree-level (Figure 4.12)

$$\frac{g_1 g_2}{\mu^2} h^\dagger D_\mu h \tilde{W} \bar{\sigma}^\mu \tilde{B}, \tag{4.12}$$

and it can also come from the dimension-five operator obtained by integrating out

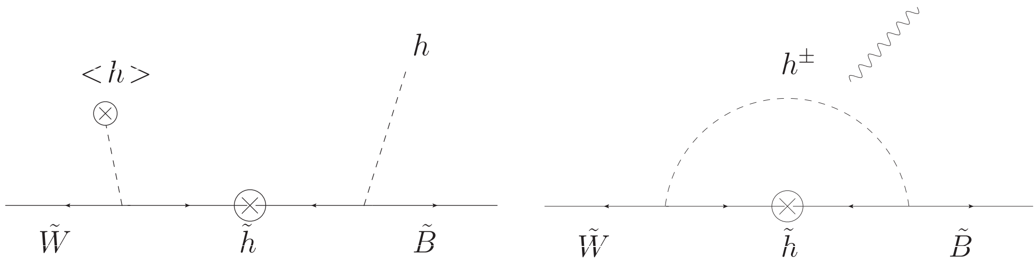

Figure 4.11: Diagrams that contribute to the dimension-five operators of Eq. (4.11) (left) and Eq. (4.13) (right).

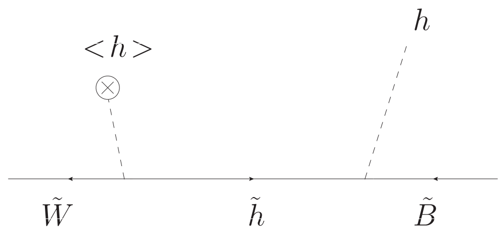

Figure 4.12: Diagram that yields the dimension-six operator of Eq. (4.12).

the Higgsinos at one-loop, generating a dipole operator (cf. right of Figure 4.11)

$$\frac{g_2^2 g_1}{16\pi^2 \mu} \tilde{W}^i_j \sigma^{\mu\nu} \tilde{B} F^j_{i\mu\nu}. \tag{4.13}$$

In either case, with heavy Higgsinos $\tilde{W}^0 \to \tilde{B}Z$ is expected to a rare decay. If the Higgsinos are heavier than ~ 10 TeV, the radiative dimension-five operator dominates the amplitude and we have a branching ratio for this decay $\sim (\alpha/(4\pi))^2 (m_{\tilde{W}}/m_Z)^2$. For heavy enough winos this could be observable. Note that the dimension-six operator can only contribute to $\tilde{W}^0 \to \tilde{B}Z$ but not to $\tilde{W}^0 \to \tilde{B}\gamma$, while the dipole operator gives both. The pure dipole predicts a ratio of the photon to Z final states of just $\sin^2\theta_W / \cos^2\theta_W \sim 1/3$. A measurement of this could establish the dipole operator as the source of the wino decay, and would show that the Higgsinos are heavy enough for the dimension-six operator to be negligible. Alternately, a deviation from this ratio would tell us that the Higgsinos are heavy, but lighter than ~ 10 TeV.

Having enumerated the decay possibilities, we now consider the signatures of gluino production and decay at the LHC. Let us assume a non-squeezed spectrum with $m_{\tilde{g}} > m_{\tilde{W}} > m_{\tilde{B}}$. This offers the possibility of spectacular processes. If the stops are the lightest colored scalars, we have the signal of $t\,\bar{t}\,t\,\bar{t}\,\tilde{B}\tilde{B}$ final states, which yields four tops + MET, where the stops are potentially produced from displaced vertices (if the scalar scale is high enough, or the spectrum is adequately squeezed). More striking is if the decay proceeds as $\tilde{g} \to t\bar{t}\tilde{W}^0$, with $\tilde{W}^0 \to \tilde{B}h$. In such a case we could find final states with 8 b's, four $W^{\pm}$ and significant MET (and again, possibly displaced vertices). Such a process would have effectively zero background, making the only question for this scenario whether gluinos are produced at all. At 14 TeV and 300 fb^{-1}, we estimate approximately 5 events for ~ 2.5 TeV gluinos (or 3 TeV gluinos for ten times that data). In some cases, the decay $\tilde{g} \to \bar{t}\,b\tilde{W}^+$ will occur, followed by $\tilde{W}^+ \to W^+\tilde{B}$. Note that this final state is very similar (topologically) to the direct decay $\tilde{g} \to \bar{t}t\tilde{B}$.

Let us now consider the possibility that the bottom of the spectrum is reversed and the wino is the LSP. Essentially all the decays should proceed via Higgs emission (if kinematically available), i.e. the decay $\tilde{g} \to \bar{t}t\tilde{B}$ will be followed by $\tilde{B} \to \tilde{W}h$. In contrast, direct decays to charged winos will proceed through $\tilde{g} \to \bar{t}\,b\tilde{W}^-$, with the chargino proceeding to decay into $\tilde{W}^0$, producing a disappearing track.

Thus, to summarize, for the $m_{\tilde{W}} > m_{\tilde{B}}$ case, the final states are $4t$+ MET, $4t\,2h$ + MET, as well as $2t\,2b\,2W$+ MET. For the $m_{\tilde{B}} > m_{\tilde{W}}$ case, the final states are $4t$ + MET, $4t\,2h$ + MET and $2t\,2b$ + MET. It is clear from this list that distinguishing these cases will be non-trivial. However, the W from the chargino decay should be distinguishable from the one that comes from top decay, while direct decay to b's should produce a spectrum of b quarks which are in principle distinct from those from top decay. Finally, the presence of the classic disappearing track signature, once seen, would be a clear sign of the wino LSP.

4.4.2 Gluino Decays and Stop Naturalness

One of the key features of an unnatural theory is that the LR soft masses (the mass mixing between a sfermion charged under $SU(2)$ and its singlet partner of the same flavor) should be negligible. Even with large A and μ, these terms are also proportional to the Higgs vev, and are thus naturally $\sim 10^4$ times smaller than the soft mass-squared terms. This impacts gluino decays in an interesting way.

In more detail, the gluino decay operators are

$$\frac{g_2}{\Lambda_L^2}\tilde{g}\, b_L\bar{t}_L\tilde{W}^- \quad \frac{g_2}{\Lambda_L^2}\tilde{g}\, t_L\bar{t}_L\tilde{W}^0 \quad \frac{g_1}{\Lambda_L^2}\tilde{g}\, t_L\bar{t}_L\tilde{B} \tag{4.14}$$

$$\frac{g_1}{\Lambda_L^2}\tilde{g}\, b_L\bar{b}_L\tilde{B} \quad \frac{g_1}{\Lambda_t^2}\tilde{g}\, t_R\bar{t}_R\tilde{B} \quad \frac{g_1}{\Lambda_b^2}\tilde{g}\, b_R\bar{b}_R\tilde{B}$$

where $\Lambda_L^{-2} = g_3 \sum_i U^L_{i3} \tilde{m}^{-2}_{L,i} U^{L*}_{i3}$, $\Lambda_t^{-2} = g_3 \sum_i U^t_{i3} \tilde{m}^{-2}_{t,i} U^{t*}_{i3}$, and $\Lambda_b^{-2} = g_3 \sum_i U^b_{i3} \tilde{m}^{-2}_{b,i} U^{b*}_{i3}$ are the mass scales obtained from a sum over squark mass eigenstates on the internal line in Figure 4.10, weighted by the matrix U that transforms from the flavor basis to the mass basis.

The key observation here is that we have five distinct decay modes into heavy flavor, $\tilde{g} \to \bar{t}t\tilde{W}^0$, $\tilde{g} \to \bar{b}b\tilde{W}^0$, $\tilde{g} \to \bar{t}b\tilde{W}^+$, $\tilde{g} \to \bar{t}t\tilde{B}$, and $\tilde{g} \to \bar{b}b\tilde{B}$. In contrast, we have only three distinct mass scales in the problem, Λ_L, Λ_t, and Λ_b. Thus, the decay of gluinos into heavy flavor is a highly overconstrained system in the unnatural limit, while for natural theories, cross terms introduce additional parameters into the theory. The heavy flavor branching ratios can easily falsify the unnatural scenario. Alternately, if they are consistent with small A terms, this would place additional fine-tuning strain on the MSSM to accommodate the Higgs mass, though of course, we cannot discount a cancellation that reduces sensitivity to these cross terms.

This discussion does not account for top polarization measurements. Should t_l and t_r be distinguishable, this would introduce yet another quantity into an already overconstrained system. If that, too, could be understood with only the three mass scales, it would provide strong evidence of a simply unnatural theory. Regardless, it would clearly show that scalar masses are not significantly corrected by electroweak symmetry breaking.

4.5 Conclusions

The expectation of a natural resolution to the hierarchy problem has always been the best reason to expect new physics at the TeV scale, accessible to the LHC. Naturalness demands new colored states lighter than a few hundred GeV, needed to stabilize the top loop corrections to the Higgs mass. These colored particles must be accessible at the TeV scale. Dark matter is another reason to expect new particles in the vicinity of the weak scale, but the WIMP "miracle" is not particularly sharp and allows for masses and cross-sections that can vary over several orders of magnitude. Indeed, if we take the simplest picture for dark matter–new electroweak doublets or triplets, annihilating through the W and Z, the needed masses are at 1 or 3 TeV, well out of range of direct production at the LHC. It is only naturalness that forces *colored* particles to be light, with the expectation that they should be copiously produced at the LHC.

On the other hand, naturalness has been under indirect pressure from the earliest days of BSM model-building, and the pressure has been continuously intensifying on a number of fronts in the intervening years. The LHC is now exploring the territory where natural new physics should have shown up. No new physics has yet been seen, and while it is far from the time to abandon the idea of a completely natural theory for electroweak symmetry breaking, the confluence of indirect and direct evidence pointing against naturalness is becoming more compelling. But the Higgs mass $m_H \sim$ 125 GeV, is within a stone's throw of its expected value in supersymmetric theories, and the compelling aspects of low-energy SUSY–precision gauge-coupling unification and dark matter–remain unaltered.

The simplest picture resolving the tensions given by this state of affairs is the minimally split SUSY model we have discussed in this chapter. These models can be easily killed experimentally, for instance, if the hint of the large enhancement to $h \to \gamma\gamma$ without enhancement to $h \to ZZ, WW$ is solidified. As for positive signals, indirect evidence for the heavy scalars can arise since since they are just in the range where they may give rise to interesting levels of FCNCs.

The direct LHC probes of these models walks on a knife's edge of excitement. Obviously if the new fermions are too heavy to be produced we have nothing. But if the gauginos are directly produced, not only do we see new particles, but since their decays can only proceed through higher-dimension operators, we get a number of handles on the presence of a high scale between 10 to 1000 TeV, ranging from displacement or flavor violation in gluino decays, to rare decay modes for the wino/bino. This would be enormously exciting, not only providing dramatic evidence for fine-tuning at the electroweak scale, but giving an indication of new thresholds that are not out of reasonable reach for future accelerators in this century.

As has long been appreciated and repeatedly pointed out, the dark matter motivation does not guarantee that the gauginos will be accessible to the LHC; the LSP could be a 3 TeV wino giving the correct relic abundance. But it is also perfectly possible that they are light enough to be produced. Fortunately, the final states from gluino decays are so spectacular that only a handful need to be produced to confirm discovery. If Nature has indeed chosen the path of un-natural simplicity, we will have to hope that she will be kind enough to let us discover this by giving us a spectrum with electroweak-inos lighter than ~ 300 GeV or gluinos lighter than ~ 3 TeV.

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Vita

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