Supersymmetry with a chargino NLSP and gravitino LSP

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ABSTRACT: We demonstrate that the lightest chargino can be lighter than the lightest neutralino in supersymmetric models with Dirac gaugino masses as well as within a curious parameter region of the MSSM. Given also a light gravitino, such as from low scale supersymmetry breaking, this mass hierarchy leads to an unusual signal where every superpartner cascades down to a chargino that decays into an on-shell W and a gravitino, possibly with a macroscopic chargino track. We clearly identify the region of parameters where this signal can occur. We find it is generic in the context of the R-symmetric supersymmetric standard model, whereas it essentially only occurs in the MSSM when $\text{sign}(M_1) \neq \text{sign}(M_2) = \text{sign}(\mu)$ and $\tan\beta$ is small. We briefly comment on the search strategies for this signal at the LHC.

KEYWORDS: Supersymmetry Phenomenology.

Contents

1. Introduction

If the fundamental scale of supersymmetry breaking is low, which can happen with gauge mediation [1], the lightest supersymmetric particle (LSP) is the gravitino. Sparticles produced at colliders will rapidly decay down to the next-to-lightest sparticle (NLSP) which then slowly decays into a particle and gravitino. This decay chain is assured assuming there is no excessively small kinematic suppression for heavier sparticles to decay into the NLSP (see e.g. [2]). The identity of the NLSP becomes paramount to determine the collider

signal; many possibilities for the NLSP have been considered [3–9] including the lightest neutralino, the stau, and the gluino.

In this paper, we demonstrate that the NLSP could be a chargino, leading to a dramatically distinct signal of supersymmetry. Every superpartner cascades down to a chargino that decays into an on-shell W and a gravitino, possibly with a macroscopic chargino track. The final decay $\tilde{\chi}^{\pm} \to W^{\pm}\tilde{G}$ is 2-body, at least for $m_{\tilde{G}} < 21\,\text{GeV}$, due to the LEPII bound on the mass of charginos, $m_{\tilde{\chi}^{\pm}} > 101\,\text{GeV}$ [10].

Common lore asserts that the lightest neutralino is always lighter than the lightest chargino in the minimal supersymmetric standard model (MSSM). This is certainly true in the bino limit $M_1 \ll M_2, \mu, M_Z$, and has been studied and confirmed in the wino limit [11–14] and Higgsino limit [15, 16], at least without excessively large radiative corrections [15]. Generically, radiative corrections to the mass *difference* between the chargino and neutralino are small, less than a GeV [17]. The exception is if the lightest gauginos are Higgsino-like with significant contributions from top and bottom squarks, where ref. [15] found it could be as large as a few GeV. It is not clear if this parameter region remains viable, in light of present direct search constraints and electroweak precision corrections. Nevertheless, as we will see, there are qualitatively distinct regions of *tree-level* gaugino parameters resulting in a chargino NLSP regardless of radiative corrections. This is the focus of the paper.

We find two qualitatively distinct scenarios where the chargino can be the NLSP. The first, and by far the most significant, is the minimal R-symmetric supersymmetric model (MRSSM). Generically, this model can have the lightest chargino lighter than the lightest neutralino due to the fundamentally different neutralino mass matrix that results from the Dirac gaugino masses. The second scenario is, remarkably, a curious and relatively unexplored region of the MSSM parameter space, where $\text{sign}(M_1) \neq \text{sign}(M_2) = \text{sign}(\mu)$ and $\tan\beta$ is small.[1] The mass difference between the lightest neutralino and the lightest chargino in the R-symmetric scenario can be tens of GeV or more, whereas it can be up to about $5\,\text{GeV}$ (at tree-level) in this curious region of the MSSM.

The organization of this paper is as follows. Section 2 reviews Dirac gauginos, the MRSSM model, and demonstrates that a chargino is NLSP in a wide region of parameter space. Section 3 is devoted to identifying the curious region of the MSSM parameter space where a chargino can be the NLSP. In section 4 the decay width of the chargino NLSP into the LSP is calculated. In section 5 the collider phenomenology of a chargino NLSP is discussed. Finally, in section 6 we conclude.

Many analytical results are presented to concretely demonstrate that the chargino can be the NLSP in the wino and Higgsino limits of the MRSSM and the MSSM. This discussion is somewhat technical; readers interested in just knowing the parameter space that results in a chargino NLSP may go directly to section 2.4 for the MRSSM (especially figures 3 and 4), and to the latter half of section 3 for the MSSM (especially figures 10

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[1]We have also found that the chargino can be the NLSP at tree-level when $\text{sign}(M_1) \neq \text{sign}(M_2) = -\mu$. In this region we find the mass difference is smaller than 1% of the NLSP mass. It is not obvious that this region remains after radiative corrections are taken into account, and thus we do not consider this region further.

and 11). Readers unfamiliar with Dirac gauginos are encouraged to read up to the end of section 2.1. Readers interested in just the new signals can skip directly to section 5.

2. Neutralinos and charginos with Dirac gaugino masses

Dirac gaugino masses result when the gaugino is married with a fermion in the adjoint representation through the operators

$$\int d^2\theta \; c_i \frac{W'_\alpha}{M} W_i^\alpha \Phi_i \;, \tag{2.1}$$

after the spurion $W'_\alpha = D\theta_\alpha$ acquires a D-term. Here Φ_i are chiral superfields in the adjoint representation of the SM groups, and c_i are dimensionless coefficients. M represents the messenger scale where these operators are generated. This possibility has been contemplated for weak scale supersymmetry some time ago [18–20] and more recently [21–28].

Gauginos which acquire Dirac masses from eq. (2.1) are not necessarily Dirac fermions once electroweak symmetry is broken and the gauginos mix with Higgsinos. Charginos are obviously Dirac fermions, since charginos carry a conserved U(1) charge, i.e., electric charge. Neutralinos are Dirac fermions only if a global U(1) is preserved by all neutralino interactions. In the minimal R-symmetric supersymmetric standard model (MRSSM) [28], a U(1)$_R$ symmetry is preserved, and thus the neutralinos are Dirac fermions. By contrast, in the Fox-Nelson-Weiner (FNW) model [21], while gauginos acquire Dirac masses, the Higgsinos acquire mass through an ordinary μ-term. The Higgsino mass violates the U(1)$_R$-symmetry, and thus leads to neutralinos that are (pseudo-Dirac) Majorana fermions. For our purposes, the most illuminating scenario with Dirac gaugino masses is the MRSSM.

The remarkable feature of the MRSSM is that it drastically ameliorates the supersymmetric flavor problem, with no excessive contributions to electric dipole moments, but with order one squark and slepton mass mixings among nearly all flavors [28, 29]. This is possible for several reasons: left-right squark and slepton mixing is absent; the gaugino masses M can be naturally $4\pi/g$ heavier than the scalar masses; and several flavor-violating operators are more suppressed than in the MSSM due to the absence of R-violating operators.

A low energy model with U(1)$_R$ symmetry can arise, for example, if supersymmetry breaking hidden sector preserves U(1)$_R$, which happens in a wide class of supersymmetry breaking models (see e.g., [30]). Nevertheless, cancellation of the cosmological constant with an R-violating constant in the superpotential [31] is generally expected to cause R-violation to be communicated from the hidden sector to the MRSSM via anomaly mediation. A natural way to minimize the size of the R-symmetry violation is to take the gravitino mass small, such as in a low scale supersymmetry breaking scenario. It is thus very natural to imagine an R-symmetric model with a light gravitino, making the resulting experimental signals important to study.

There are several other models that have Dirac gaugino masses where we find that the chargino NLSP phenomenon can also occur. In appendix A we briefly comment on the FNW model [21], showing that there are specific limits where the neutralinos become Dirac fermions and the mass matrices reduce to the ones found in the MRSSM. Hence, our results apply to this model as well.

Fields	$S(3)_C$	$S(2)_W$	$U(1)_Y$	$U(1)_R$
Q	3	2	$\frac{1}{6}$	1
U	$\bar{3}$	1	$-\frac{2}{3}$	1
D	$\bar{3}$	1	$\frac{1}{3}$	1
L	1	2	$-\frac{1}{2}$	1
E	1	1	1	1
$\Phi_{\tilde{B}}$	1	1	0	0
$\Phi_{\tilde{W}}$	1	3	0	0
$\Phi_{\tilde{g}}$	8	1	0	0
H_u	1	2	$\frac{1}{2}$	0
H_d	1	2	$-\frac{1}{2}$	0
R_u	1	2	$-\frac{1}{2}$	2
R_d	1	2	$\frac{1}{2}$	2

Table 1: Gauge and R-charges of all chiral supermultiplets in the MRSSM.

2.1 Mass matrices of the MRSSM

In the MRSSM, gaugino masses arise from eq. (2.1) which generate Dirac masses that pair the gauginos $(\tilde{g}, \tilde{W}$ and $\tilde{B})$ with their fermionic partners $(\psi_{\tilde{g}}, \psi_{\tilde{W}}$ and $\psi_{\tilde{B}})$. Higgsino masses arise from pairing Higgs superfields H_u and H_d with partner fields R_u and R_d through a pair of mass terms

$$\int d^2\theta \left[\mu_u H_u R_u + \mu_d H_d R_d \right] . \tag{2.2}$$

$R_{d,u}$ transform identically to $H_{u,d}$ under the electroweak group, except that the R-charges are 2 rather than 0. This R-charge assignment forbids Yukawa-like couplings of the R-fields to the matter fields. Hence, it is natural to assume that only Higgses acquire electroweak symmetry breaking expectation values. Upon electroweak symmetry breaking, the electroweak gauginos mix with the Higgsinos (just like in the MSSM). For completeness, all the multiplets in the MRSSM described in ref. [28] are listed in table 1 along with their matter and R-charges.

Let us first investigate the neutralino mass matrix in the MRSSM. The R-charges determine which neutralinos mix with each other and provide a guiding principle to determine the gauge-eigenstate basis. The vector $N_+ \equiv \left(\tilde{W}_3, \tilde{B}, \tilde{R}_u^0, \tilde{R}_d^0 \right)$ carries R-charge $+1$, while the vector $N_- \equiv \left(\psi_{\tilde{W}}^0, \psi_{\tilde{B}}, \tilde{H}_u^0, \tilde{H}_d^0 \right)$ carries a R-charge -1. Field rotations do not mix up fields with different R-charges. Hence, just like with charginos in the MSSM, two independent rotation matrices are required to diagonalize the mass matrix.

The Lagrangian for neutralino masses can be written in the gauge-eigenstate basis as

$$\mathcal{L}_{\text{neutralino mass}} = N_+^{\text{T}} M_{\tilde{N}} N_- + \text{c.c.} , \tag{2.3}$$

where the neutralino mass matrix is given as

$$
M_{\tilde{N}} = \begin{bmatrix}
M_2 & 0 & -gv_u/\sqrt{2} & gv_d/\sqrt{2} \\
0 & M_1 & g'v_u/\sqrt{2} & -g'v_d/\sqrt{2} \\
-\lambda_u v_u/\sqrt{2} & \lambda'_u v_u/\sqrt{2} & \mu_u & 0 \\
\lambda_d v_d/\sqrt{2} & -\lambda'_d v_d/\sqrt{2} & 0 & \mu_d
\end{bmatrix} .
\tag{2.4}
$$

Here $\langle H_u \rangle \equiv v_u$ and $\langle H_d \rangle \equiv v_d$ with $v_u^2 + v_d^2 = v^2/2 \simeq (174 \text{ GeV})^2$. Notice the apparently unusual location of the μ-terms in eq. (2.4) is a direct result of the Dirac nature of the neutralino mass matrix. The physical mass-squareds are given by eigenvalues of $M_{\tilde{N}} M_{\tilde{N}}^{\mathrm{T}}$. There are four new parameters $\lambda_u, \lambda_d, \lambda'_u$ and λ'_d that arise from the superpotential terms

$$
\int d^2\theta \left[H_u \left(\lambda_u \Phi_{\tilde{W}} + \lambda'_u \Phi_{\tilde{B}} \right) R_u + H_d \left(\lambda_d \Phi_{\tilde{W}} + \lambda'_d \Phi_{\tilde{B}} \right) R_d \right] .
\tag{2.5}
$$

These couplings are unnecessary to the structure of the MRSSM but are nevertheless allowed under all of the charge assignments. The λ couplings also open up the possibility of spontaneously breaking R-symmetry through renormalization. In this paper we consider a low energy effective theory that preserves R-symmetry. Various checks have been performed on eq. (2.4) to verify that every entry in this mass matrix is correct, see appendix B.

For charginos, the mass matrix is even simpler because of the conservation of electromagnetic charge as well as R-charge. In the MRSSM there are eight two-component fermions from the winos, Higgsinos, and R-fields. Based on the R-charges and the electromagnetic charges, the eight fermions can be grouped into following four different classes which do not mix with each other:

charges	$Q = +1$	$Q = -1$
$R = +1$	$\chi_{++} \equiv \left(\tilde{W}^+, \tilde{R}_d^+ \right)$	$\chi_{+-} \equiv \left(\tilde{W}^-, \tilde{R}_u^- \right)$
$R = -1$	$\chi_{-+} \equiv \left(\psi_{\tilde{W}}^+, \tilde{H}_u^+ \right)$	$\chi_{--} \equiv \left(\psi_{\tilde{W}}^-, \tilde{H}_d^- \right)$

$$\tag{2.6}$$

These charges imply that the charginos pair up as

$$
\mathcal{L}_{\text{chargino mass}} = \chi_{++}^{\mathrm{T}} M_{\chi 1} \chi_{--} + \chi_{-+}^{\mathrm{T}} M_{\chi 2} \chi_{+-} + \text{c.c.} ,
\tag{2.7}
$$

where the chargino mass matrices in our basis are

$$
M_{\chi 1} = \begin{bmatrix} M_2 & gv_d \\ \lambda_d v_d & \mu_d \end{bmatrix} \quad \text{and} \quad M_{\chi 2} = \begin{bmatrix} M_2 & \lambda_u v_u \\ gv_u & \mu_u \end{bmatrix} .
\tag{2.8}
$$

Since each of these matrices are diagonalized by bi-unitary transformations, four independent rotations are needed for the four vectors listed in eq. (2.6) in order to diagonalize the mass matrices.

We are now ready to calculate the mass eigenstates under various assumptions about the parameters of the MRSSM. One qualitative difference from the MSSM is that the Dirac nature of the gauginos allows us to take the gaugino masses M_1, M_2 and Higgsino masses

μ_u, μ_d real and positive (rotating phases into the holomorphic masses for the scalar adjoints and the λ parameters) [28]. Examining the gaugino mass matrices, we explore several limits where we can analytically demonstrate that the chargino is the NLSP and obtain a good estimate of the mass difference. Our first treatment is to take $\lambda_{u,d} = \lambda'_{u,d} = 0$. This is motivated in part to simplify our analysis, but also to emphasize that nonzero values of these couplings are *not* necessary to obtain a chargino lighter than the lightest neutralino. Further simplifications can result by taking the large $\tan\beta$ limit and the large M_1 limit. In section 2.3 we generalize and expand the discussion, while still working at tree-level. As we will see, the lightest chargino can be *significantly* lighter than the lightest neutralino when the λ couplings are present with $\mathcal{O}(g)$ values.

2.2 Simplified MRSSM: large M_1 with $\lambda = \lambda' = 0$

The gaugino mass matrices, eqs. (2.4) and (2.8), simplify in the limit $\lambda_{u,d} = \lambda'_{u,d} = 0$. Here we will also take $\mu_u = \mu_d = \mu$, which will prove extremely convenient in our analytic analysis in this section. Equal Higgsino masses will also allow us to illustrate the contrast between the MSSM and the MRSSM. We are obviously not interested in the case where the bino is the lightest gaugino, hence we take large M_1, consistent with the motivations of ref. [28]. Integrating out the bino and taking $\tan\beta \gg 1$, the neutralino and chargino mass matrices are given by

$$
M_{\tilde{N}} = \begin{bmatrix} M_2 & -M_W & 0 \\ 0 & \mu & 0 \\ 0 & 0 & \mu \end{bmatrix} \qquad \text{and} \qquad M_{\chi 2} = \begin{bmatrix} M_2 & \sqrt{2}M_W \\ 0 & \mu \end{bmatrix} . \tag{2.9}
$$

For charginos, in the case $\tan\beta > 1$, we need only consider $M_{\chi 2}$ in order to find the lightest chargino mass.

The block diagonal form of $M_{\tilde{N}}$ in eq. (2.9) shows that a pair of neutral Higgsinos acquire a (Dirac) mass μ and do not mix with the other neutralinos. This allows us to further simplify the neutralino mass matrix down to just the upper 2×2 block. Here is the key analytical observation: The upper 2×2 block of $M_{\tilde{N}}$ is identical to $M_{\chi 2}$, except that the off-diagonal element is *smaller* for the neutralino mass matrix. Simple 2×2 diagonalization clearly shows that the lightest chargino is lighter than the lightest neutralino in *both* the Higgsino limit ($\mu < M_1, M_2$) and the wino limit ($M_2 < \mu, M_1$),

$$
\begin{aligned}
\text{Higgsino limit:} \quad \Delta_+ &= -\frac{\mu M_W^2}{2M_2^2} + \mathcal{O}\left(\frac{1}{M_2^4}\right) , \\
\text{wino limit:} \quad \Delta_+ &= -\frac{M_2 M_W^2}{2\mu^2} + \mathcal{O}\left(\frac{1}{\mu^4}\right) .
\end{aligned} \tag{2.10}
$$

where $\Delta_+ \equiv m_{\tilde{\chi}_1^\pm} - m_{\tilde{\chi}_1^0}$.

It is well known that the one-loop electromagnetic radiative correction increases the chargino mass [11–14]. But, since the size of the mass difference shown above is not parametrically small compared to the loop contribution, the tree-level splitting can easily dominate so long as one of the diagonal elements does not far exceed the other.
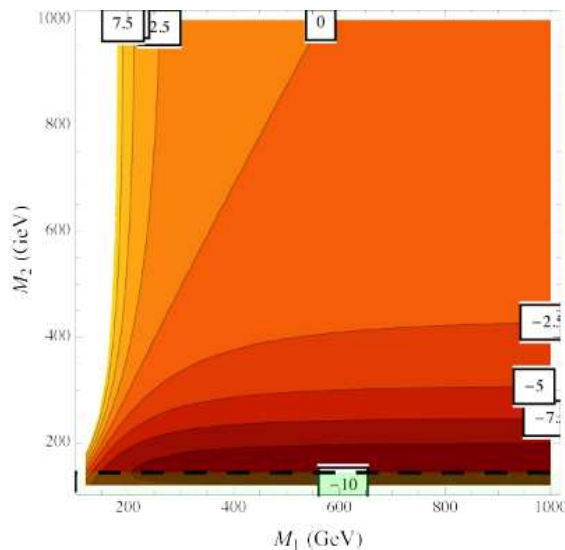
Figure 1: Contours of Δ_+ (GeV) in the simplified MRSSM at $\tan\beta = 10$ and $\mu = 150\,\text{GeV}$.

The analytical results are strikingly confirmed by examining a larger region of the MRSSM parameter space numerically. In figures 1 and 2 we plot the contours of Δ_+, the difference of the lightest chargino mass to the lightest neutralino mass. Figure 1 explores the mostly-Higgsino limit and was generated by holding $\mu = 150\,\text{GeV}$ and $\tan\beta = 10$. Similarly, figure 2 explores the mostly-wino limit and was generated holding $M_2 = 150\,\text{GeV}$ and $\tan\beta = 10$. The regions under the dashed lines in these figures result in $m_{\tilde{\chi}_1^\pm} < 101\,\text{GeV}$ at tree level and thus will be ignored from further consideration.

In these figures, it is clear that a chargino is the lightest gaugino in the regions with $\Delta_+ < 0$. This occurs throughout the tree-level parameter space of the wino limit, shown in figure 2. Note that we have not included radiative corrections in these numerical results because a full calculation requires knowing the full spectrum of the model. Nevertheless, it is clear that in a wide range of parameter space, the tree-level mass difference is much larger than are expected from radiative corrections, demonstrating that the chargino can indeed be the NLSP.

2.3 Enlarging the parameter space: nonzero λ couplings

The full parameter space of the MRSSM is much larger than the MSSM. To keep things manageable we consider the following simplifications: $\lambda_u = \lambda_d \equiv \lambda$ and $\lambda_u' = \lambda_d' \equiv \lambda'$. Moreover, it will be useful to define the parameters $M_\lambda = \lambda v/2$ and $\tan\theta_\lambda = \lambda'/\lambda$, analogous to the electroweak parameters $M_W = gv/2$ and $\tan\theta_W = g'/g$. We use perturbation techniques on this mass-squared matrix to find approximate analytical expressions for neutralino masses.
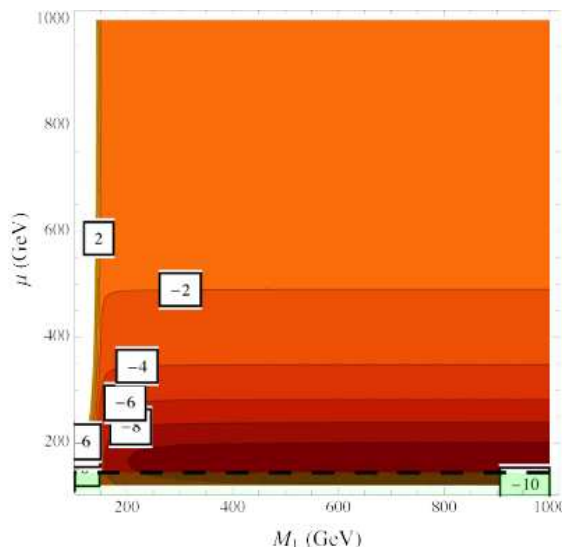
Figure 2: Contours of $\Delta_+(\mathrm{GeV})$ in the simplified MRSSM at $\tan\beta = 10$ and $M_2 = 150\,\mathrm{GeV}$.

2.3.1 Equal higgsino masses $\mu_u = \mu_d = \mu$

In this case, the neutralino mass matrix simplifies in a very interesting way:

$$
M_{\tilde{N}} = \begin{bmatrix}
M_2 & 0 & -\sin\beta\,M_W & \cos\beta\,M_W \\
0 & M_1 & \sin\beta\,\tan\theta_W\,M_W & -\cos\beta\,\tan\theta_W\,M_W \\
-\sin\beta\,M_\lambda & \sin\beta\,\tan\theta_\lambda\,M_\lambda & \mu & 0 \\
\cos\beta\,M_\lambda & -\cos\beta\,\tan\theta_\lambda\,M_\lambda & 0 & \mu
\end{bmatrix} . \quad (2.11)
$$

The lower 2×2 block is proportional to the identity matrix, which implies that an orthogonal transformation involving only the last two columns and rows will leave the lower 2×2 block invariant. Moreover, the ratio of the third and the fourth elements in the first two rows and as well as in the first two columns is $\tan\beta$. Therefore a rotation of the last two rows and columns by an angle β will make the third element of the first two rows and columns vanish simultaneously while leaving the lower 2×2 block unchanged. In this new basis, we then have zeroes everywhere in the third row and column except at the diagonal position. Thus we have managed to decouple a pair of neutralinos of mass μ from the other eigenstates.

The chargino mass-squared are similarly found after diagonalizing $M_{\chi_1}M_{\chi_1}^{\mathrm{T}}$ and $M_{\chi_2}M_{\chi_2}^{\mathrm{T}}$. These rank two matrices are straightforward to evaluate. For the purpose of comparing with the neutralino masses, however, we expand the eigenvalues in various limits. When $\tan\beta > 1$, the lightest chargino mass is found from M_{χ_2}. For the mass difference, we find

- **Higgsino limit:** $M_1, M_2 > \mu$

$$
\Delta_+ = -\,M_\lambda M_W \left(\frac{2\sin^2\beta - 1}{M_2} - \frac{\tan\theta_W\tan\theta_\lambda}{M_1} \right) + \mathcal{O}\left(\frac{1}{M_2^2}, \frac{1}{M_1^2} \right) , \quad (2.12)
$$

- **wino limit:** $M_1, \mu > M_2$

$$\Delta_+ = -\left(2\sin^2\beta - 1\right)\frac{M_\lambda M_W}{\mu} + \mathcal{O}\left(\frac{1}{\mu^2}, \frac{1}{M_1^2}\right) . \tag{2.13}$$

To leading order in this expansion, a chargino is clearly the lightest gaugino in the wino limit. In the Higgsino limit, the ratio of M_2 and M_1 is important. When $M_1/M_2 > \tan\theta_W \tan\theta_\lambda/(2\sin^2\beta - 1)$, we again find a lighter chargino.

2.3.2 One higgsino heavy: $\mu_d \gg M_1, M_2$ and $\mu_u = \mu$

Another limit which can be analyzed analytically is when one Higgsino is much heavier than the other mass parameters. Taking $\mu_d \gg M_1, M_2$, we can immediately integrate out down-type Higgsinos resulting in a 3×3 neutralino matrix. For charginos, we need to take into account only M_{χ_2}. Again using perturbation techniques, we find

- **Higgsino limit:** $M_1, M_2 > \mu$

$$\Delta_+ = -M_\lambda M_W \sin^2\beta \left(\frac{1}{M_2} - \frac{\tan\theta_W \tan\theta_\lambda}{M_1}\right) + \mathcal{O}\left(\frac{1}{M_2^2}, \frac{1}{M_1^2}\right), \tag{2.14}$$

- **wino limit:** $M_1, \mu > M_2$

$$\Delta_+ = -\sin^2\beta \frac{M_\lambda M_W}{\mu} + \mathcal{O}\left(\frac{1}{\mu^2}, \frac{1}{M_1^2}\right) . \tag{2.15}$$

The mass differences calculated in this case are quite similar to the case of equal Higgsino masses, except for the dependency on $\sin^2\beta$. In fact, in the large $\tan\beta$ limit, eqs. (2.14) and (2.15) reduce to eqs. (2.12) and (2.13) respectively. A careful gaze reveals that, at sizeable $\tan\beta$, the relevant portion of the gaugino mass matrices are identical whether one considers the equal Higgsino case or one Higgsino heavy case.

2.4 Numerical results

Having demonstrated analytically that the chargino can be the NLSP in several limits, we now turn to analyzing a larger region of the MRSSM parameter space numerically. The neutralino and chargino masses are determined by nine parameters: M_1, M_2, μ_u, μ_d, $\tan\beta$, λ_u, λ_d, λ'_u, and λ'_d. As it is too cumbersome to do a complete scan, we restrict to the simplifications introduced in the previous subsection: $\lambda_u = \lambda_d = \lambda$ and $\lambda'_u = \lambda'_d = \lambda'$. As before, we also trade the parameters λ and λ' for the mass parameters M_λ and the angle θ_λ.

In the first part of this subsection we keep M_λ and $\tan\theta_\lambda$ fixed and scan the rest of the parameter space in order to understand the variation of Δ_+. Later we will choose a particular point in the M_1, M_2, μ_u, μ_d and $\tan\beta$ space, and see the dependence of Δ_+ on our choice of M_λ and $\tan\theta_\lambda$. To begin our discussion, we take $\tan\theta_\lambda = \tan\theta_W$ and $M_\lambda = M_W$.
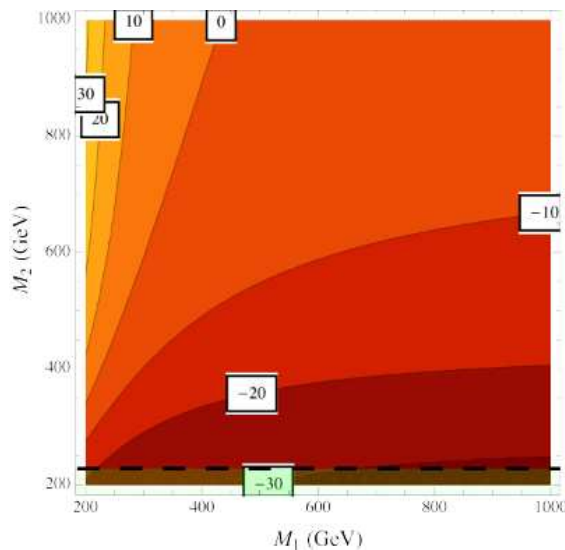
Figure 3: Contours of Δ_+(GeV) in the MRSSM at $\tan\beta = 10$ and $\mu = 200\,\mathrm{GeV}$ in the equal Higgsino mass limit.
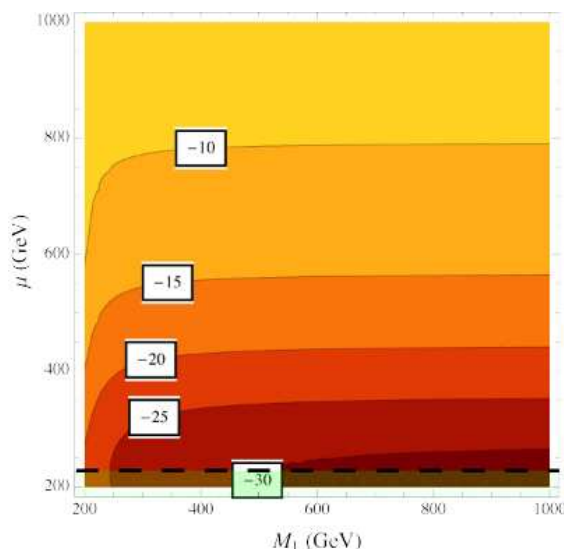

Figure 4: Contours of Δ_+(GeV) in the MRSSM at $\tan\beta = 10$ and $M_2 = 200\,\mathrm{GeV}$ in the equal Higgsino mass limit.

First, consider the case of equal Higgsino masses with $\tan\beta = 10$. (The one heavy Higgsino scenario, as we described before, is similar to the equal Higgsino scenario at sizeable $\tan\beta$). In figure 3 we hold μ constant and vary M_2 and M_1. The shape of the contours are the same as in figure 1, where we neglected the "λ" couplings. The striking difference is that now the mass difference (Δ_+) can be as large as $-30\,\mathrm{GeV}$. In fact, this is why we chose $\mu = 200\,\mathrm{GeV}$, so that the lightest chargino remains above the LEPII limit. In figure 4, the same analysis is repeated in the wino limit, with $M_2 = 200\,\mathrm{GeV}$, where again we see that mass difference Δ_+ is negative and can be up to $-30\,\mathrm{GeV}$.
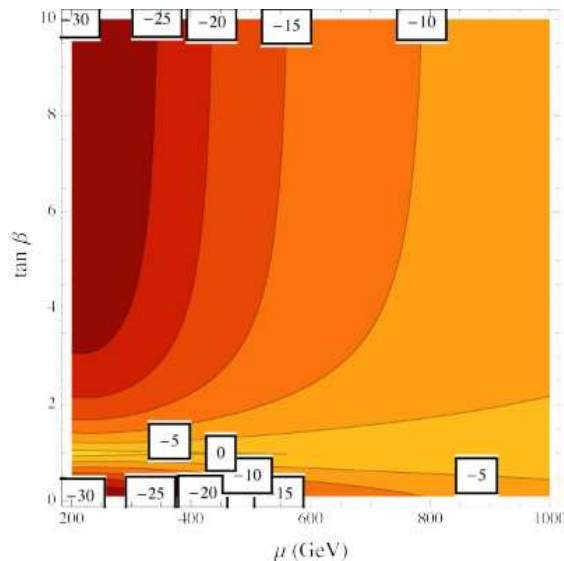
Figure 5: Contours of Δ_+(GeV) in the MRSSM at $M_1 = 500\,\text{GeV}$ and $M_2 = 200\,\text{GeV}$ in the equal Higgsino limit.
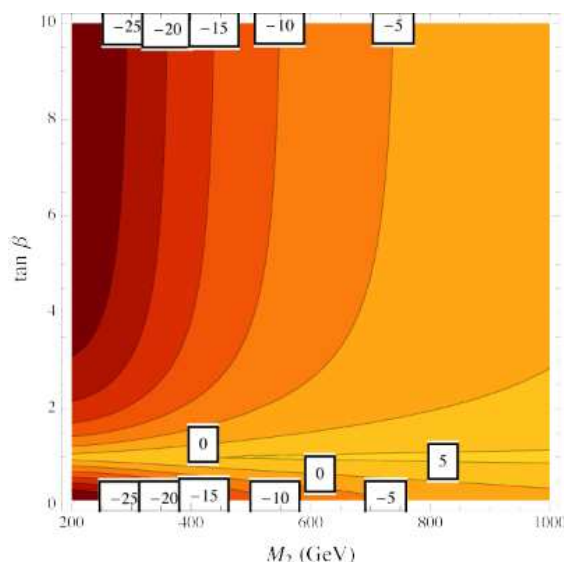
Figure 6: Contours of Δ_+(GeV) in the MRSSM at $M_1 = 500\,\text{GeV}$ and $\mu = 200\,\text{GeV}$ in the equal Higgsino limit.

It is interesting to investigate the dependence of Δ_+ on $\tan\beta$ in the wino and Higgsino limits. In figure 5, we take $M_1 = 500\,\text{GeV}$, $M_2 = 200\,\text{GeV}$, and vary μ and $\tan\beta$. In figure 6, we take $M_1 = 500\,\text{GeV}$, $\mu = 200\,\text{GeV}$, and vary M_2 and $\tan\beta$. Both the figures show similar features. For a given set of mass parameters, the contours are largely insensitive to the value of $\tan\beta$ as long as $\tan\beta$ is sizeable. Close to $\tan\beta = 1$ the contours change rapidly. At $\tan\beta = 1$ we find $\Delta_+ \geq 0$ in both the wino and Higgsino cases. The case $\tan\beta < 1$ is symmetric with respect to $\tan\beta > 1$ upon taking $\tan\beta \to 1/\tan\beta$.
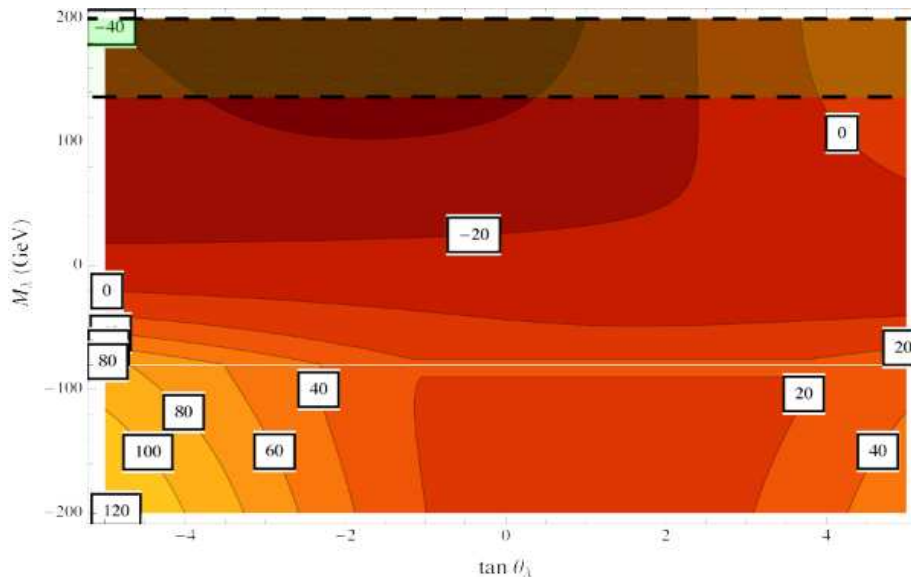
Figure 7: Contours of Δ_+(GeV) in the $\tan\theta_\lambda - M_\lambda$ plane for $M_1 = 600$ GeV, $\tan\beta = 10$ and $\mu = M_2 = 250$ GeV.

We can also investigate how M_λ and $\tan\theta_\lambda$ affects the chargino/neutralino mass hierarchy. This is most easily done by considering a specific point in the M_1, M_2, μ and $\tan\beta$ space where Δ_+ is large. From figures 3 and 4, we find large (negative) Δ_+ when $M_2 \sim \mu$. In addition, Δ_+ is almost independent of M_1 for large enough M_1. Finally, from figures 5 and 6 we see that large Δ_+ occurs at large $\tan\beta$.

In figure 7, the variation of Δ_+ in the $M_\lambda - \tan\theta_\lambda$ plane is shown. We take $M_1 = 600\,\text{GeV}$ and $\tan\beta = 10$ with $\mu = M_2 = 250$ GeV. We do not vary M_λ beyond ± 200 GeV, since even at this value λ already exceeds 1. The figure clearly shows that we get the largest (negative) splitting when M_λ is positive and $\tan\theta_\lambda$ is negative, up to $\Delta_+ \sim -\mathcal{O}(50\,\text{GeV})$. The shaded area above the dashed line is excluded since it results in a chargino with mass $m_{\tilde\chi_1^\pm} < 101\,\text{GeV}$. If $\mu = M_2$ is reduced, Δ_+ becomes larger (negative). On the other hand, the upper bound on M_λ from the bound on the chargino mass also decreases rapidly. For larger values of $\mu = M_2$, the upper bound on M_λ increases and effectively one can also find a larger splitting.

2.5 Summary and comments

From various analytical results and numerical figures it is clear that a chargino can be significantly lighter than the lightest neutralino. The mass difference can far exceed the size of radiative corrections to the gaugino masses.

- The expressions for the mass differences between the lightest neutralino and the lightest chargino in eq. (2.12) and in eq. (2.13) were obtained for $\tan\beta > 1$. The case $\tan\beta < 1$ can be obtained by substituting $\sin\beta \leftrightarrow \cos\beta$.

- The expressions for Δ_+ are given in eq. (2.14) and in eq. (2.15) when μ_d is large. An alternate limit where μ_u is greater than all other relevant masses may be found by substituting $\sin\beta \leftrightarrow \cos\beta$ in eq. (2.14) and in eq. (2.15).

- Given the results from the various limits, combined with figures 5 and 6, we find that at $\tan\beta = 1$, a neutralino is always the lightest gaugino.

- Finally, in the absence of λ couplings, the mass difference Δ_+ can be obtained in the Higgsino limit *without* decoupling M_1. Starting with eqs. (2.12) and (2.14), hold $\tan\theta_\lambda$ finite, while independently taking $M_\lambda \to 0$. We obtain

$$\mu_u = \mu_d = \mu \quad \to \quad \Delta_+ = -\,\mu\, M_W^2 \left(\frac{2\sin^2\beta - 1}{M_2^2} - \frac{\tan^2\theta_W}{M_1^2} \right) ,$$
$$\text{Large } \mu_d \quad \to \quad \Delta_+ = -\,\sin^2\beta\,\mu\, M_W^2 \left(\frac{1}{M_2^2} - \frac{\tan^2\theta_W}{M_1^2} \right) . \tag{2.16}$$

These expressions demonstrate the chargino can be lighter than the neutralino in the MRSSM without λ couplings.

3. Neutralino and chargino masses in the MSSM

We now turn to studying the neutralino and chargino masses in the MSSM. The neutralino mass matrix in the MSSM is rank four, and although exact analytical expressions for the eigenvalues exist [32–34], they are not particularly transparent. We instead consider several well-known limits where the mass difference between the chargino and neutralino can be easily calculated analytically. Later in this section we generalize our results using numerical calculations. We allow the mass parameters M_1, M_2, μ to have arbitrary sign, though without loss of generality we can take $M_2 > 0$. We do not consider arbitrary phases, since they are severely constrained in the MSSM from the absence of electric dipole moments [35]. The neutralino gauge eigenstates include a bino with mass M_1, hence in the small M_1 limit a mostly-bino neutralino will always be the lightest gaugino. The nontrivial cases of interest to us will occur when M_1 is not the smallest parameter in the mass matrix.

The two interesting limits that could have a chargino NLSP are the Higgsino limit and the wino limit. In the Higgsino limit $M_1, M_2 > \mu, M_W$, calculations of the masses of light gauginos including radiative corrections can be found in [15, 16]. The mass difference between the lightest chargino and lightest neutralino is

$$\Delta_+ = \left[\left(\tan^2\theta_W \frac{M_2}{M_1} + 1 \right) + \left(\tan^2\theta_W \frac{M_2}{M_1} - 1 \right) \frac{\mu}{|\mu|} \sin 2\beta \right] \frac{M_W^2}{2M_2} + \mathcal{O}\left(\frac{1}{M_2^2} \right) , \tag{3.1}$$

The neutralino-chargino mass splitting grows as M_2 is reduced.

In the wino limit $M_2 < M_1, \mu$, the lightest neutralino and the lightest chargino are highly degenerate. The mass difference was calculated including one-loop effects in [11–14]. The tree-level splitting can be obtained by expanding in powers of $1/\mu$, with the leading order splitting occurring at $\mathcal{O}\left(1/\mu^2\right)$:

$$\Delta_+ = \frac{M_W^2}{\mu^2} \frac{M_W^2}{M_1 - M_2} \tan^2\theta_W \sin^2 2\beta + \mathcal{O}\left(\frac{1}{\mu^3} \right) . \tag{3.2}$$
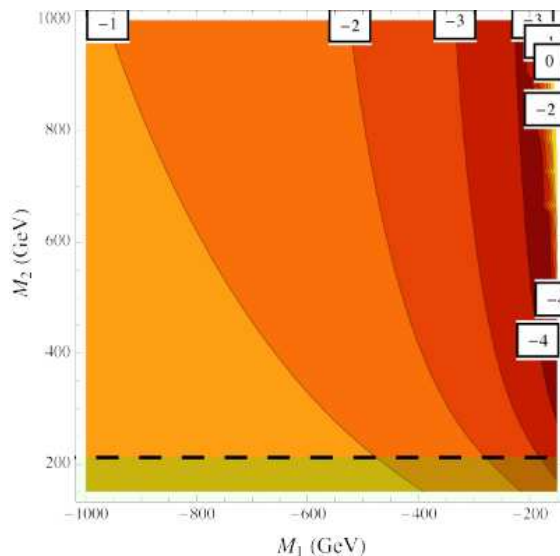
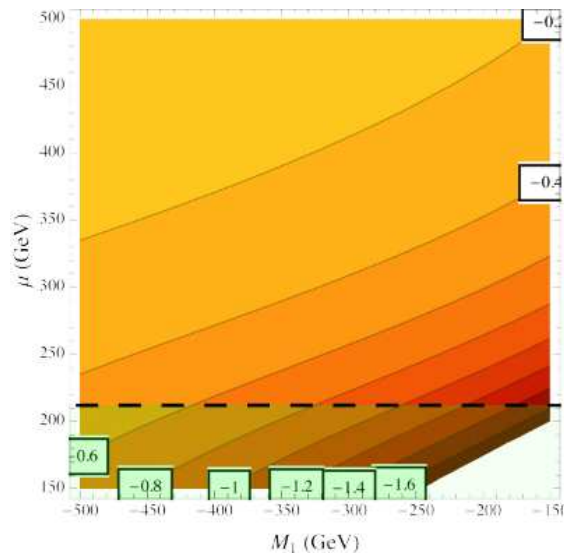**Figure 8:** Contours of $\Delta_+(\text{GeV})$ in the MSSM at $\tan\beta = 2$ and $\mu = 150\,\text{GeV}$.

Figure 9: Contours of $\Delta_+(\text{GeV})$ in the MSSM at $\tan\beta = 2$ and $M_2 = 150\,\text{GeV}$.

One loop corrections to eq. (3.2) are positive and typically small, of order $0.1\,\text{GeV}$ [14].

In the well-known case $\text{sign}(M_1) = \text{sign}(M_2)$, it is evident from the leading terms in eqs. (3.1) and (3.2) that $m_{\tilde{\chi}_1^\pm} > m_{\tilde{\chi}_1^0}$ for any sign of μ in either the Higgsino limit or wino limit. The lightest neutralino persists as the lightest gaugino throughout the MSSM parameter space with $\text{sign}(M_1) = \text{sign}(M_2)$. For completeness and comparison to the MRSSM, we demonstrate this explicitly in appendix C.

More interestingly, when $M_1 < 0$ and $M_2 > 0$, the tree-level mass difference between the lightest chargino and the lightest neutralino is no longer positive definite. This is clear already at leading order in both the Higgsino limit eq. (3.1) and the wino limit eq. (3.2).
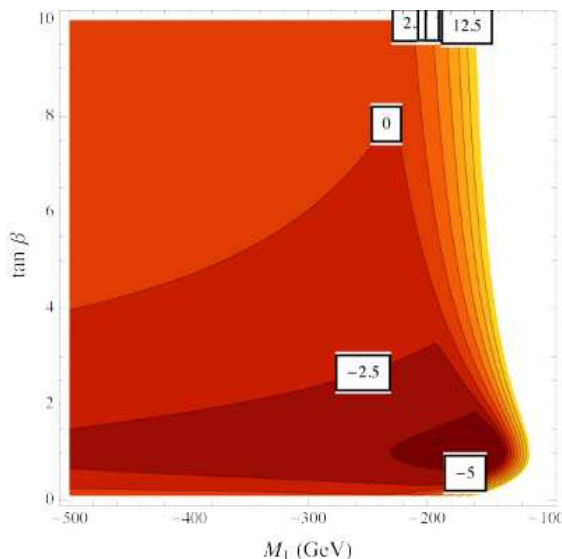
Figure 10: Contours of Δ_+(GeV) in the MSSM at $M_2 = 600\,\text{GeV}$ and $\mu = 150\,\text{GeV}$.

The mass difference Δ_+ grows as $\tan\beta \to 1$, and it is also generally larger in the Higgsino limit. To calculate size of the splitting, especially in regions of the parameter space not covered by the limits given above, we turn to numerical evaluation of the masses.

The following figures summarize the numerical results. First we fix $\tan\beta = 2$. This is about the smallest value of $\tan\beta$ that could possibly yield a lightest Higgs boson mass above the LEP II bound [36]. In figures 8 and 9 we take $\mu = 150\,\text{GeV}$ and $M_2 = 150\,\text{GeV}$ respectively. The regions under the dotted lines in these figures result in $m_{\tilde\chi_1^\pm} < 101\,\text{GeV}$ at tree level and we ignore them from further consideration. When μ in figure 8 and M_2 in figure 9 are reduced further, the dashed line for $M_{\tilde\chi_1^\pm} = 101\,\text{GeV}$ rises.

As before, we show contours of Δ_+. In both of these figures there are considerable regions with negative Δ_+, demonstrating that the chargino can be NLSP. Notice that the splitting is largest in the Higgsino limit, up to about $5\,\text{GeV}$. We deliberately show a smaller range of values of M_1 and μ in figure 9 so that the gradient of Δ_+ may be clearly seen. Here M_2 is held fixed and we start with large μ. Δ_+ increases as μ is decreased. The rate of increase of splitting peaks as μ becomes comparable to M_2.

We show the $\tan\beta$ dependence of the size of the splitting in figure 10. We set $M_2 = 600\,\text{GeV}$ and $\mu = 150\,\text{GeV}$, which is a point in figure 8 where the splitting is the largest. We find that the splitting between lightest neutralino and the lightest chargino is maximized at $\tan\beta = 1$ for a fixed value of M_1 and for $M_1 \simeq -\mu$ for a given $\tan\beta$. For the sake of completeness we also evaluated the mass splitting in the wino limit, for the parameters $M_2 = 150\,\text{GeV}$ and $\mu = 250\,\text{GeV}$, in figure 11. Not surprisingly, the splitting is smaller than that in the Higgsino limit. However, the purpose of the figure is to demonstrate that a wino-like chargino NLSP also occurs at small $\tan\beta$.
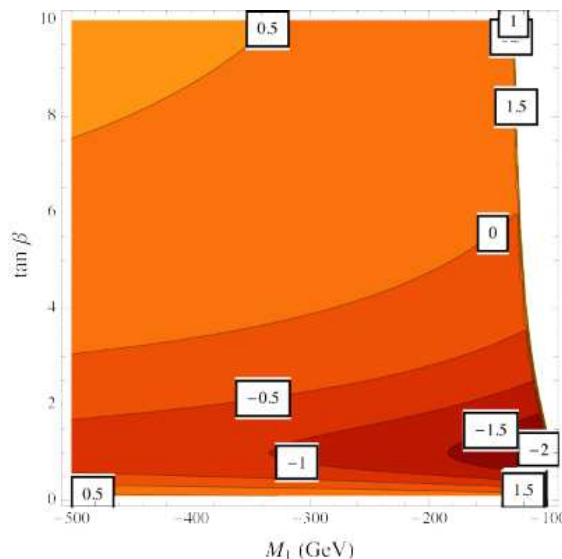
Figure 11: Contours of $\Delta_+(\text{GeV})$ in the MSSM at $M_2 = 150\,\text{GeV}$ and $\mu = 250\,\text{GeV}$.

Finally, we comment on the gaugino mass hierarchy within a common extension of the MSSM that incorporates a gauge singlet S. In the next-to-minimal supersymmetric standard model (NMSSM), the superpotential includes

$$W \supset \lambda_S S H_u H_d + \frac{1}{6} k S^3 \,, \tag{3.3}$$

where the Higgsino mass $\mu = \lambda_S \langle S \rangle$ is generated after S acquires an scalar expectation value. The chargino mass matrix is unchanged while the neutralino mass matrix is enlarged to a 5×5 matrix containing additional parameters: the singlino mass $m_S = k\mu/\lambda_S$ and the singlino-Higgsino mixing λ_S. Of these parameters, λ_S can always be chosen to be positive, while the sign of m_S is arbitrary (as is the sign of μ). Like the MSSM, we find the neutralino is heavier than the chargino whenever all mass parameters (M_1, M_2 and m_S) are positive. However, if $m_S < 0$ the chargino can be the NLSP even if $M_1, M_2 > 0$. This result can be understood by observing that bino mixing is entirely analogous to singlino mixing with other neutralino states, replacing the bino mass with the singlino mass and the $g'v_{u,d}$ Higgsino mixing terms by $\lambda_S v_{u,d}$. Therefore, for the purposes of determining eigenvalues, the singlino acts just like the bino. As a result, the chargino can be lightest when either M_1 or m_S is negative. However, just like the MSSM, the chargino-neutralino splitting remains small, $\lesssim 5$ GeV.

4. Chargino decay

Having demonstrated that the chargino can be the NLSP, we now turn to considering the decay of the chargino into the gravitino. Given that the mass splitting between the chargino NLSP and a neutralino next-to-next-to-lightest supersymmetric particle (NNLSP) can be small, it may be possible for the neutralino NNLSP to decay directly to a gravitino. Hence,

in the discussion below, we consider both chargino and neutralino 2-body decays to the gravitino. In the next section we will compare the rates for 2-body decay of the NNLSP directly into a gravitino against the 3-body decay of the NNLSP into the NLSP.

Charginos could decay to $W^{\pm}$ plus gravitino, or to an electrically charged (± 1, R-charge neutral) scalar plus gravitino. Similarly, neutralinos decay into a photon, Z^0 or a neutral scalar plus gravitino. In the MSSM, these scalars are contained in H_u, H_d, while in the MRSSM, there are extra scalars in $\Phi_{\tilde{W}}, \Phi_{\tilde{B}}$. In addition R_u, R_d contain scalars, but they carry R-charge 2 and cannot be involved in the decay. The form of the decay width of a chargino (or neutralino) $\tilde{\chi}$ into a gravitino and a spectator particle (X) is [5]:

$$\Gamma(\tilde{\chi} \to \tilde{G} + X) \sim \frac{\kappa \, m_{\tilde{\chi}}^5}{96\pi M_{pl}^2 \tilde{m}_{3/2}^2} \left(1 - \frac{m_X^2}{m_{\tilde{\chi}}^2} \right)^4 \tag{4.1}$$

where κ is an order one mixing angle. The decay width (4.1) is sensitive to the mass of the spectator particle (m_X), and decays to heavier final states are kinematically suppressed. The kinematic factor makes the decays to the lightest particle possible, $W^{\pm}$ in the case of a chargino and γ for a neutralino, the preferred mode. Additionally, the charged scalar mass matrices contains at least one extra parameter, namely b (or m_{A^0}), compared to the gaugino mass matrices. By adjusting the additional parameter(s) we are always free to focus on the simpler scenario where decays to charged scalars are kinematically forbidden. Converting the above width to a decay length for a chargino at rest,

$$L = 3.4 \, \kappa^{-1} \left(\frac{100 \text{ GeV}}{m_{\tilde{\chi}}} \right)^5 \left(\frac{\tilde{m}_{3/2}}{10 \text{ eV}} \right)^2 \left(1 - 0.6467 \left(\frac{m_{\tilde{\chi}}}{100 \text{ GeV}} \right)^{-2} \right)^{-4} \text{ mm} \tag{4.2}$$

The characteristic decay length is proportional to $m_{3/2}^2$. In weakly-coupled messenger sectors, the gravitino mass can be related to the sparticle masses. If the mediation scale is low, the gravitino is typically too light to produce a visible charged track. With strong interactions present in the hidden sector, however, the gravitino mass can be significantly enhanced with respect to the rest of the spectrum [37–39], opening up the possibility of a chargino track. In the following, the gravitino mass is taken to be a free parameter, and thus our analysis applies regardless of the scale of messenger interactions.

Two important applications of formula (4.1) for us are: the decay of the lightest chargino through $\tilde{\chi}^{\pm} \to W^{\pm}\tilde{G}$, and the decay of the lightest neutralino into $\gamma +$ gravitino, $Z+$ gravitino. Within the MRSSM, working in the Higgsino limit to leading order in $1/M_2$:

$$\kappa_{W\tilde{G}} = \sin^2\beta, \quad \kappa_{\gamma\tilde{G}} = \sin^2\theta_W \sin^2\beta \, \frac{M_W^2}{M_2^2}, \quad \kappa_{Z\tilde{G}} = \sin^2\beta \, . \tag{4.3}$$

Alternatively, in the wino limit we have:

$$\kappa_{W\tilde{G}} = 1, \qquad \kappa_{\gamma\tilde{G}} = \sin^2\theta_W, \qquad \kappa_{Z\tilde{G}} = \cos^2\theta_W \, . \tag{4.4}$$

to leading order in $\mathcal{O}(1/\mu)$. To lowest order in either limit, the κ couplings are independent of λ, λ'.

When M_2 is small, the lightest neutralino is mostly the neutral wino, and thus the decays to γ, Z are simply proportional to the photino or zino fraction of $\tilde{W}_3$. On the other hand, when μ is small and the lightest neutralino is primarily a Higgsino, the coupling to the Z remains order one, while the coupling to the photon is suppressed.

Shifting to the MSSM, in the limit $\mu \sim -M_1$, and $|\mu|, |M_1| \ll |M_2|$ the lightest neutralino is a maximal mixture of Higgsino and bino, while the chargino is purely Higgsino. This content is reflected in the κ mixing angles:

$$\kappa_{W\tilde{G}} = 1 + \sin^2 2\beta \, \frac{M_W^2}{M_2^2} \ ,$$

$$\kappa_{\gamma\tilde{G}} = \cos^2\theta_W \left[1 + \frac{1}{4\sqrt{2}} \tan\theta_W \, (1 - \sin 2\beta) \, \frac{M_W}{M_1} \right] \ ,$$

$$\kappa_{Z\tilde{G}} = \frac{1}{2} \, \cos^2\beta \, + \, \sin^2\theta_W \left[1 \, + \, \frac{1}{4\sqrt{2}} \tan\theta_W \Big(\frac{1 - \sin 2\beta}{\cos\beta + \sin\beta} \Big) \frac{M_W}{M_1} \right]$$
$$- \, \frac{1}{16\sqrt{2}} \tan\theta_W \left(\frac{2\cos\beta - 5\sin\beta + 3\sin 3\beta}{1 \, + \, \tan\beta} \right) \frac{M_W}{M_1} \ .$$

5. Collider phenomenology

In low-scale supersymmetry models, the NLSP is typically either a neutralino or a charged slepton. As all particles in a low-energy supersymmetry-breaking model eventually decay down to the NLSP, its properties such as spin, mass, charge, and decay width form the foundation upon which all collider studies are built. By demonstrating that a chargino can be the NLSPs, we are opening the door to an entirely new class of sparticle signatures, with a plethora of exciting phenomenological consequences.

Rather than study a particular exclusive process, we focus here on the inclusive cross sections for sparticle production with chargino NLSP. A detailed study of the optimal cuts to pick out a given sparticle spectrum over the background is beyond the scope of this paper; instead, our aim is to identify search channels and the most important SM backgrounds. This effort is in the same spirit as ref. [5], which studied the inclusive signal with a neutralino NLSP, namely $\gamma\gamma + \not{E}_T$. We also comment on the possibilities for distinguishing the chargino NLSP from other potential NLSPs.

5.1 Inclusive sparticle signal

We model the inclusive sparticle signal by the total production cross section for all twelve species of squarks. We do not consider the leptons or Higgses, and assume all squarks to be degenerate in mass. The inclusive squark production cross section as a function of the common squark mass is shown in figure 12 below.

In this scenario, the squarks first decay into a quark plus chargino or neutralino. The subsequent decays of the heavy (non-NLSP) charginos/neutralinos depend on the details of the gaugino spectrum: if kinematically allowed, heavy gauginos decay to a light gaugino plus a gauge bosons, otherwise they will decay into three-body decay final state containing a light gauginos plus two fermions. Finally, the chargino NLSPs each decay
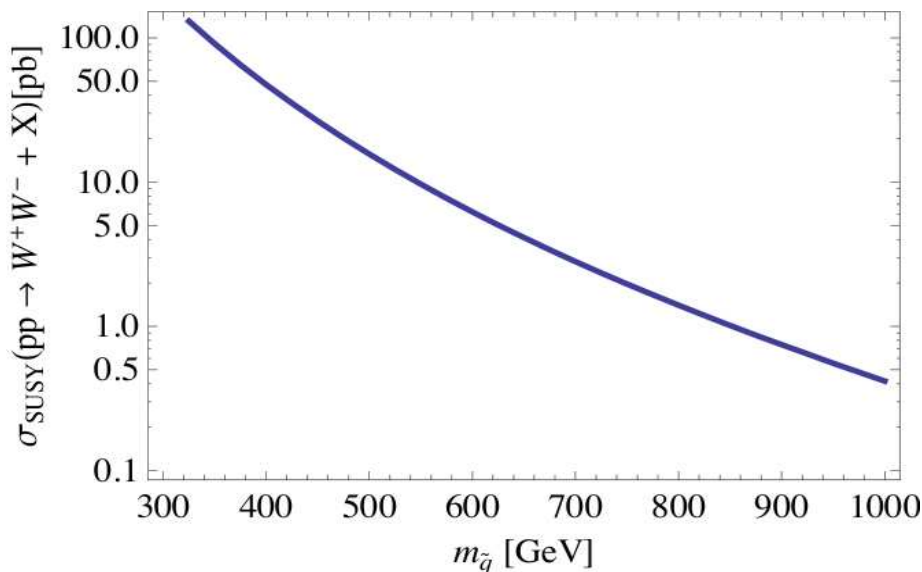
Figure 12: Lowest order cross section $\sigma_{\text{SUSY}}(pp \to W^+W^- + X)$ at the LHC (14 TeV). We have approximated the inclusive sparticle cross section by the QCD production cross section of all 12 squarks. For simplicity the squarks were taken to be degenerate, and sleptons were assumed to be heavier than the squarks.
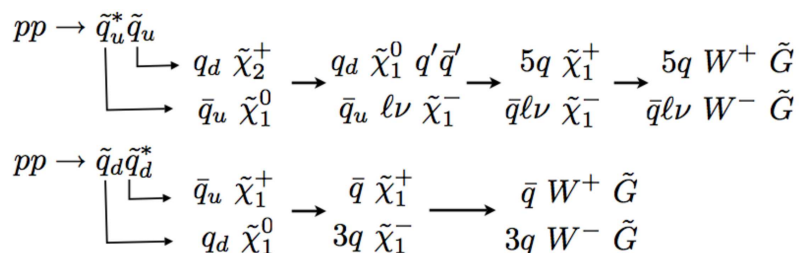
$$
\begin{array}{l}
pp \to \tilde{q}_u^* \tilde{q}_u \\
\qquad \begin{array}{l}
\;\vdash\!\!\to\; q_d\; \tilde{\chi}_2^+ \to q_d\; \tilde{\chi}_1^0\; q'\bar{q}' \to 5q\; \tilde{\chi}_1^+ \to 5q\; W^+\; \tilde{G} \\
\;\vdash\!\!\to\; \bar{q}_u\; \tilde{\chi}_1^0 \quad\;\; \bar{q}_u\; \ell\nu\; \tilde{\chi}_1^- \quad\; \bar{q}\ell\nu\; \tilde{\chi}_1^- \quad\; \bar{q}\ell\nu\; W^-\; \tilde{G}
\end{array} \\[2ex]
pp \to \tilde{q}_d \tilde{q}_d^* \\
\qquad \begin{array}{l}
\;\vdash\!\!\to\; \bar{q}_u\; \tilde{\chi}_1^+ \to \bar{q}\; \tilde{\chi}_1^+ \longrightarrow \bar{q}\; W^+\; \tilde{G} \\
\;\vdash\!\!\to\; q_d\; \tilde{\chi}_1^0 \quad\; 3q\; \tilde{\chi}_1^- \qquad 3q\; W^-\; \tilde{G}
\end{array}
\end{array}
$$

Figure 13: Possible decay chains resulting from squark production. Exactly what path is typical depends on details of the sparticle spectrum.

into W + gravitino. Thus, *every* supersymmetric event contains at least two on-shell Ws plus missing energy. Using this model for sparticle production, we want to discuss the sparticle discovery potential via excesses in $W^+W^- + \not{E}_T + X$. Some possible decay chains are shown in figure 13.

Where $\tilde{\chi}_2^+, \tilde{\chi}^0$ are the second-lightest chargino and lightest neutralino respectively. As we can see, in addition to on-shell W pairs, inclusive sparticle events involve cascades containing additional (hard) quarks and leptons. These extra leptons/jets are especially important in scenarios where the chargino lives longer than ~ 25 ns: by the time the long-lived charginos decay to objects which can be triggered on, the next bunch-crossing will have occurred in the detector, causing the charginos to be associated with the wrong event. However, if there are extra leptons/jets present in the sparticle decays, one can trigger on those objects instead of on the chargino decay products. Then, by refining the analysis offline to search for charged tracks, the presence of charginos could be revealed.

Model	Limit	$\xi_L^2 + \xi_R^2$
MRSSM	$\mu \ll M_2, M_1$	$\frac{1}{2}$
MRSSM	$M_2 \ll \mu, M_1$	2
MSSM	$\mu \sim -M_1, M_1 \ll M_2$	$\frac{1}{4}\left[1 - \frac{1}{4\sqrt{2}}\tan\theta_W \left(\frac{1-\sin 2\beta}{\sin\beta+\cos\beta}\right)\frac{M_W}{M_1}\right]$

Table 2: The parameter values in different limits and models that determine the ratio R_Γ.

5.2 NNLSP decay

Another complication in this scenario is the possibility that the second-heaviest chargino and/or the lightest neutralino decay directly into gravitino $+ X$ rather than decay via three-body decays. If the heavier chargino decays directly, the signal will still contain two W bosons plus a pair a gravitinos, while if a neutralino decays directly one of the Ws is replaced by a hard photon or Z.[2] The effect of the direct gravitino decays depends on the relative rates:

$$
\begin{aligned}
R_\Gamma = \frac{\Gamma(\tilde{\chi}_H \to \tilde{\chi}_L f f')}{\Gamma(\tilde{\chi}_H \to X + \tilde{G})} &= N_f \frac{4g^4 M_{pl}^2 \tilde{m}_{3/2}^2}{5\pi^2 M_W^4}\left(\frac{\xi_L^2 + \xi_R^2}{\kappa}\right)\left(\frac{\Delta}{m_{\tilde\chi}}\right)^5 \\
&= 0.556 N_f \left(\frac{\xi_L^2 + \xi_R^2}{\kappa}\right)\left(\frac{\tilde{m}_{3/2}}{10\text{ eV}}\right)^2 \left(\frac{100\text{ GeV}}{m_{\tilde\chi}}\right)^5 \left(\frac{\Delta}{2\text{ GeV}}\right)^5,
\end{aligned}
\tag{5.1}
$$

where Δ is the mass difference between the heavy chargino/neutralino ($\tilde{\chi}_H$) and the NLSP ($\tilde{\chi}_L$): $\Delta = m_{\tilde{\chi}_H} - m_{\tilde{\chi}_L}$. The prefactor N_f is the number of fermionic degrees of freedom $\tilde{\chi}_L$ is kinematically allowed to decay into, and $\xi_L, \xi_R, (\kappa)$ are mixing angles from the chargino-neutralino-X ($\tilde{\chi}_H$-gravitino-X) interactions; values in the limits of interest are given in table 2. We plot R_Γ for the neutralino as a function of the mass splitting Δ and gravitino mass in figure 14 below.

The region where direct decays are important is restricted to light gravitinos and small splittings. For gravitinos heavier than a few tenths of an eV, the three body decays are always dominant in either of the MRSSM limits we consider. For the MSSM, the large bino component of the lightest neutralino leads to a large $\kappa_{\gamma\tilde{G}}$ and suppressed $\xi_{L,R}$, thus direct decays can be important for gravitinos as heavy as 1 eV. Neglecting direct decays of heavy gauginos to gravitinos, we now explore various search strategies and backgrounds to $W^+W^- + \not{E}_T$.

5.3 Search strategies and backgrounds

The experimental search strategy will depend greatly on the lifetime of the chargino. If the chargino is long-lived but decays within the detector, it will leave a charged track, possibly leading to a displaced vertex. As mentioned above, triggering on the long-lived

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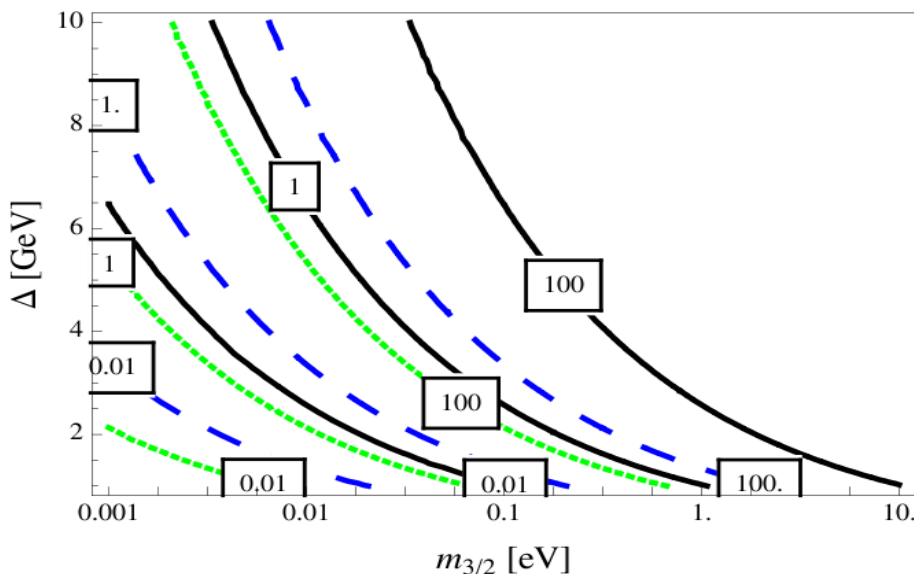[2]Neutral Higgses are also a possibility, which we ignore for simplicity here.

Figure 14: Decay ratio R_Γ as a function of the gravitino mass and the mass splitting Δ taking the NLSP mass to be $100\,\mathrm{GeV}$. $R_\Gamma < 1$ corresponds to NNLSPs dominantly decaying directly into a gravitino rather than through the NLSP. The dotted line (green) refers to the MRSSM Higgsino limit; the dashed line (blue) refers to the MRSSM wino limit; and the solid line (black) refers to the MSSM with $M_1 = -200$ GeV, $\tan\beta = 2$). These three cases correspond to the three limiting cases of table 2. In all three cases the contours indicate values, from left to right, of $R_\Gamma = 0.01$, $R_\Gamma = 1$, and $R_\Gamma = 100$.

charginos is troublesome. However, provided they can be found, displaced vertices are great discriminant of new physics from background.[3]

If the charginos decay promptly, conventional variables such as $H_T = \sum_{i=\ell,j} p_T$ and $M_{\mathrm{eff}} = H_T + \not{E}_T$ [40] will likely be the first to indicate discovery. Additionally, these variables will be effective handles for separating signal from background. The optimum value for cuts in these variables depends on the superpartner mass scale.

In the case of prompt charginos, the dominant SM backgrounds in $W^+W^-\not{E}_T + X$ are: $\bar{t}t + \mathrm{jets}$, single top + jets, $W^+W^-/W^\pm Z + \mathrm{jets}$, and $W^\pm/Z + \mathrm{jets}$.[4] Exactly which background is most important will depend on the strategy for W pair detection.

Irreducible SM backgrounds that result from $W^+W^- \to \ell + \nu + jj$, contain only one true source of missing energy - the neutrino from the leptonic W decay. In addition to providing less $\not{E}_T$ to SM events (compared to two neutrinos), with a single source of missing energy one can reconstruct the transverse mass of the W ($m_{T,W}$) from the p_T of the lepton and the $\not{E}_T$. For the SM events the transverse mass distribution will exhibit a peak near M_W, followed by steep dropoff. Meanwhile, the m_T distribution for the signal will fall much slower (after M_W) due to the excess $\not{E}_T$ carried by the gravitinos. Therefore, by cutting on

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[3]Lifetimes in the vicinity of 5 ns would actually be ideal: 5 ns is small enough that triggering/resolution is not an issue, yet long enough that timing information in the calorimeter could be used to differentiate signal from background. We thank Dirk Zerwas for bringing this point to our attention.

[4]By "+jets" we are including heavy quark flavors (b, c) as well as light.

$m_{T,W} \gtrsim 100$ GeV we can remove a large fraction of the background while maintaining the signal. The signal may still be polluted by events where two leptons are produced but one of them is missed by the detector. An important background which falls into this category is $\bar{t}t \to \ell\nu\ell'\nu'\bar{b}b$.[5]

A second channel to search for W pairs, $\ell\nu\ell'\nu'$, has the advantage of an additional lepton, which greatly reduces the $W^{\pm}/Z +$ jets background. Unfortunately, when both W's decay leptonically, the SM backgrounds have higher missing energy, so the efficiency of $\not{E}_T$ cuts will be reduced. The p_T and η of the two leptons remain useful variables for discriminating signal from background.

The third possibility is for both Ws to decay hadronically. In that case, the signal is $4+$ jets $+ \not{E}_T$, with four of the jets breaking up into two pairs, each pair reconstructing to a W. Although reconstructing Ws can be difficult with realistic jet resolution, the background can be greatly suppressed by imposing hard cuts (100 GeV or more) on the p_T of the reconstructed Ws and the missing energy. However, unless there is an additional lepton in these events from an earlier cascade decay, multijet QCD becomes an important and difficult background. Further detailed study, in all of the channels mentioned above as well as for a more diverse spectrum are needed.

5.4 Confounding NLSPs

If the chargino NLSP is long-lived, a charged track will be visible in the detector. Depending on the exact lifetime of the chargino, the charged track will either exit the detector or end in a displaced vertex. Either way, a visible track allows us to rule out the majority of neutralino NLSP scenarios. However, distinguishing a chargino NLSP from a charged slepton NLSP requires more careful study.

Slepton decay proceeds through a 2-body channel to $l\tilde{G}$, whereas chargino decays through $W\tilde{G}$ with the confounding channel $W \to l\nu$. This kinematic difference may provide a useful discriminant especially when the NLSP has a larger boost. Additionally, W bosons decay to jets and (democratically) into all three leptons - therefore chargino decays will lead to e, μ, τ each in equal amounts, up to reconstruction and detector efficiency effects. Selectron and smuon NLSPs also decay to lepton plus $\not{E}_T$, but their final state is a particular lepton flavor (for at least approximately flavor-diagonal slepton masses). By simply counting the number and type of leptons in a data sample, one should easily be able to distinguish selectrons/smuons from charginos. Stau NLSPs are somewhat trickier because taus decay democratically to electrons and muons; however, provided adequate tau-tagging capability at the LHC, staus should also be easily distinguishable from charginos.

Short-lived charginos are somewhat trickier. In principle, the identity of the particle produced along with the gravitino can tell us something about the NLSP; for neutralinos this spectator particle is likely a photon, for a chargino it will be a W, and for sleptons the spectator is a lepton. Additionally, the maximum energy achieved by the spectator depends on the LSP mass, allowing one to distinguish WIMP LSP scenarios from gravitino

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[5]The $\bar{t}t$ and single-top backgrounds can be reduced further by rejecting any events with a b-tag, at the expense of compromising the stop/sbottom contribution to (12), while $W/Z +$ jets can be reduced by requiring two jets to reconstruct a W.

LSP scenarios [4]. In practice, the success of spectator-identification will depend on the details of the sparticle spectrum, and it is easy to dream up tricky scenarios.

To distinguish the chargino NLSP limit of the (N)MSSM from the MRSSM additional observables are needed. A discriminant is the mass difference between the lightest neutralino and the lightest chargino. In the MSSM the NLSP chargino occurs when several gauginos are nearly degenerate, thus it is difficult to arrange mass differences (even with large radiative corrections) greater than ~ 10 GeV. In the MRSSM large neutralino-chargino mass differences are easier to accommodate, especially if $\lambda, \lambda' \neq 0$.

6. Conclusions

In this paper we have studied a new signal of supersymmetry that results when a chargino is the NLSP and the gravitino is LSP. A necessary condition for this scenario to occur is that a chargino must be the lightest gaugino. We have found

- The chargino can be the lightest gaugino in a wide range of parameter space when neutralinos are Dirac fermions, such as in the MRSSM.

- In the MSSM a chargino can be the NLSP essentially only in the case $\text{sign}(M_1) \neq \text{sign}(M_2) = \text{sign}(\mu)$.

- There is qualitative difference between the generated gaugino mass hierarchies depending on whether the neutralinos are Dirac fermions or Majorana fermions. In the MSSM, the splitting is large when $\tan\beta$ is small, as opposed to the case of Dirac gauginos when the splitting is maximized for large $\tan\beta$.

- In addition we also observe quantitative difference between the two cases. In the MSSM the splitting is small ($\lesssim \mathcal{O}(5 \text{ GeV})$) and in the MRSSM the splitting can be as big as $\mathcal{O}(30 \text{ GeV})$ at tree level.

Given a light gravitino, a chargino NLSP will decay into a gravitino and an on-shell W. Summarizing the phenomenology:

- If the gravitino mass is far larger than $\mathcal{O}(100 \text{ eV})$, a chargino produced at a collider will almost always escape the detector leaving a charged track. (Rarely, the chargino can decay anywhere within the detector.)

- If the gravitino mass is of order $\mathcal{O}(1 - 100 \text{ eV})$, a chargino produced at a collider can have a displaced vertex and/or track, resulting in a decay well away from the interaction vertex.

- If the gravitino mass is far smaller than $\mathcal{O}(1 \text{ eV})$, a chargino produced at a collider will decay promptly into a W and a gravitino.

In the first two cases, the charged track provides a great way to discriminate signal from background. Distinguishing a chargino NLSP from a slepton NLSP requires exploiting the smaller boost and flavor-democratic decay of the W. When the chargino decays promptly,

the pair of Ws in each event provide a striking signature of sparticle production. Standard methods may suffice to reduce background and extract the signal, however additional benefit could be achieved by taking advantage of the characteristic $WW + \not{E}_T + X$ final state. Dedicated studies of exclusive sparticle production may provide more promising opportunities to discover supersymmetry with a chargino NLSP.

Acknowledgments

A. The Fox-Nelson-Weiner model with Dirac gaugino masses

The FNW model [21] contains the same content as the MRSSM *without* the R_u, R_d fields. Gaugino masses arise exclusively from the "supersoft" operator, eq. (2.1). The Higgs sector in this model, however, has an ordinary μ-term in place of eq. (2.2). Just like the MSSM, the μ-term marries the Higgsinos $\tilde{H}_u$ and $\tilde{H}_d$ with each other, and after expanding around the nonzero vevs of the two Higgses, one finds mass terms between the gauginos and the Higgsinos. The resulting neutralino mass terms are given by

$$\mathcal{L} \supset \frac{1}{2} N_0^{\mathrm{T}} M_n N_0 \,, \tag{A.1}$$

where

$$N_0 = \begin{bmatrix} \tilde{W}_3 \\ \tilde{B} \\ \tilde{H}_d^0 \\ \psi_{\tilde{W}}^0 \\ \psi_{\tilde{B}} \\ \tilde{H}_u^0 \end{bmatrix} \quad \text{and} \quad M_n = \begin{bmatrix} 0 & 0 & gv_d/\sqrt{2} & M_2 & 0 & -gv_u/\sqrt{2} \\ 0 & 0 & -g'v_d/\sqrt{2} & 0 & M_1 & g'v_u/\sqrt{2} \\ gv_d/\sqrt{2} & -g'v_d/\sqrt{2} & 0 & 0 & 0 & \mu \\ M_2 & 0 & 0 & 0 & 0 & 0 \\ 0 & M_1 & 0 & 0 & 0 & 0 \\ -gv_u/\sqrt{2} & g'v_u/\sqrt{2} & \mu & 0 & 0 & 0 \end{bmatrix} . \tag{A.2}$$

The fields ψ_i are the fermions in the supermultiplets Φ_i. In particular, $\psi_{\tilde{W}}$ is a S(2)$_W$ triplet containing a neutral $\psi_{\tilde{W}}^0$ and as well as charged components $\psi_{\tilde{W}}^\pm \equiv \frac{1}{\sqrt{2}}\left(\psi_{\tilde{W}}^{(1)} \pm i\psi_{\tilde{W}}^{(2)}\right)$. The chargino mass terms in this model are given by:

$$\mathcal{L} \supset \begin{bmatrix} \tilde{W}^+ & \psi_{\tilde{W}}^+ & \tilde{H}_u^+ \end{bmatrix} \begin{bmatrix} M_2 & 0 & gv_d \\ 0 & M_2 & 0 \\ 0 & gv_u & \mu \end{bmatrix} \begin{bmatrix} \psi_{\tilde{W}}^- \\ \tilde{W}^- \\ \tilde{H}_d^- \end{bmatrix} . \tag{A.3}$$

Similar to the discussion of the MRSSM, in the limit of large $\tan\beta$ (where $v_d \to 0$) the neutralino mass matrix simplifies drastically. A linear combination of $U(1)_R$ under which the gauginos are charged and $U(1)_Y$ under which the Higgsinos are charged is preserved and thus neutralinos become Dirac fermions with the following mass matrix:

$$\mathcal{L} \supset \left[\tilde{W}_3 \ \tilde{B} \ \tilde{H}_d^0 \right] \begin{bmatrix} M_2 & 0 & -gv_u/\sqrt{2} \\ 0 & M_1 & g'v_u/\sqrt{2} \\ 0 & 0 & \mu \end{bmatrix} \begin{bmatrix} \psi_{\tilde{W}}^0 \\ \psi_{\tilde{B}} \\ \tilde{H}_u^0 \end{bmatrix} . \tag{A.4}$$

The chargino mass matrix also simplifies:

$$\mathcal{L} \supset \left[\tilde{W}^+ \ \psi_{\tilde{W}}^+ \ \tilde{H}_u^+ \right] \begin{bmatrix} M_2 & 0 & 0 \\ 0 & M_2 & 0 \\ 0 & gv_u & \mu \end{bmatrix} \begin{bmatrix} \psi_{\tilde{W}}^- \\ \tilde{W}^- \\ \tilde{H}_d^- \end{bmatrix} . \tag{A.5}$$

Despite $U(1)_R$ broken by the Higgs sector in the model, it is easy to see that a chargino can be the NLSP. In the large $\tan\beta$ limit, the mass matrices in eqs. (A.4) and (A.5) are identical to what was obtained in the MRSSM in the case when the two Higgsinos have equal masses (i.e. $\mu_u = \mu_d = \mu$). Similarly, when $\mu_d \gg M_1, M_2, \mu_u = \mu$ and $\tan\beta$ is large, the relevant part of the mass matrices in the MRSSM in eqs. (2.4) and (2.8) become identical to eqs. (A.4) and (A.5) respectively.

B. Verifying the form of the MRSSM gaugino mass matrices

One important verification results by setting all supersymmetry breaking parameters as well as all λ parameters to zero. In this limit, the masses of the neutralinos and charginos arise due to the vev of the Higgses and their kinetic terms. The kinetic terms of the Higgses are obviously independent of R-symmetry. This implies that the masses of the heavy charginos and neutralinos should be identical whether in the MSSM or in the MRSSM. This requirement verifies that the elements in the mass matrices proportional to gauge couplings have the correct form.

In order to verify the rest of the couplings we expanded the corresponding terms in the Lagrangian in terms of electromagnetic eigenstates in the place of weak eigenstates. The numerical factors between the couplings of the charge neutral states and the charged states now determine various elements in the mass matrices.

$$\begin{aligned} H_u R_u \quad &\rightarrow \quad H_u^0 R_u^0 \ + \ H_u^+ R_u^- \\ H_u \Phi_{\tilde{W}} R_u \quad &\subset \quad \frac{1}{\sqrt{2}} H_u^0 \left(\frac{1}{\sqrt{2}} \Phi_{\tilde{W}}^0 R_u^0 \ + \ \Phi_{\tilde{W}}^+ R_u^- \right) \\ W \Phi_{\tilde{W}} \quad &\rightarrow \quad W^+ \Phi_{\tilde{W}}^- + W^- \Phi_{\tilde{W}}^+ + W_3 \Phi_{\tilde{W}}^0 , \end{aligned} \tag{B.1}$$

where the charge eigenstates are defined the usual way.

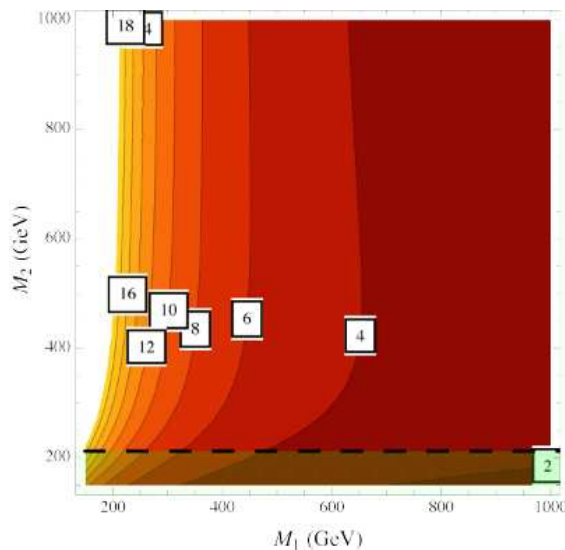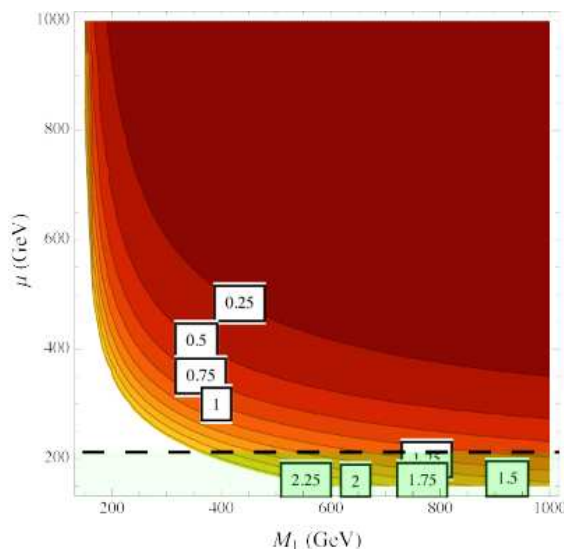**Figure 15:** Contours of Δ_+ (GeV) in the MSSM at $\tan\beta = 2$ and $\mu = 150\,\mathrm{GeV}$.

Figure 16: Contours of Δ_+ (GeV) in the MSSM at $\tan\beta = 2$ and $M_2 = 150\,\mathrm{GeV}$.

C. Chargino and neutralino masses in the MSSM with $\mathrm{sign}(M_1) = \mathrm{sign}(M_2)$

It is illuminating to numerically explore a broader range of the MSSM parameter space. In figures 15 and 16 we plot the contours of Δ_+, the difference of the lightest chargino mass to the lightest neutralino mass. Figure 15 explores the mostly-Higgsino limit and was generated by holding $\mu = 150\,\mathrm{GeV}$ and $\tan\beta = 2$. Similarly, figure 16 explores the mostly-wino limit and was generated holding $M_2 = 150\,\mathrm{GeV}$ and $\tan\beta = 2$. In each case, the figures clearly show $\Delta_+ > 0$ throughout the parameter space, and thus the neutralino is always lighter than the chargino. This confirms the results of the limiting cases eq. (3.1) and

eq. (3.2). Moreover, these figures also demonstrate that a neutralino remains the lightest gaugino even when all the parameters in the mass matrices are of the same order. The regions under the dashed lines in these figures result in $m_{\tilde{\chi}_1^\pm} < 101\,\text{GeV}$ at tree level. When μ in figure 15 and M_2 in figure 16 are reduced further, the dashed line for $M_{\tilde{\chi}_1^\pm} = 101\,\text{GeV}$ rises. For example setting $\mu = 120\,\text{GeV}$ excludes the entire region below $M_2 = 450\,\text{GeV}$.

The numerical results carry additional insight into the chargino/neutralino system. A common feature of figure 15 and figure 16 is that the contours decrease as one moves to larger values of both the $x-$ and $y-$coordinates (to the top-right corner). In figure 15 this can be understood by noting that if both M_1, M_2 are heavy, the gauginos associated with them (the bino and winos) can be integrated out; the low energy effective theory is left with just Higgsinos, which contain a pair of degenerate neutralinos and charginos, i.e. $\Delta_+ \to 0$ in the decoupling limit. Exactly the same arguments go through for the wino limit (with increasing μ and M_1), shown in figure 16 where the decoupling is much more rapid. This provides additional verification of our analytical results in eq. (3.1) where we saw that the splitting was more suppressed in the wino case.

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