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A minimal supersymmetric scenario with only μ at the weak scale

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Abstract: Inspired by split supersymmetry, we study a minimal supersymmetric scenario with only the Higgsino mass parameter μ below the TeV scale. The motivation is to satisfy the gauge coupling unification and dark matter constraints with the minimal particle contents at the electroweak scale in supersymmetric models. With the neutral Higgsino as the lightest supersymmetric particle, we discuss the dark matter signals in both direct and indirect detection. We also discuss collider phenomenology associated with the two lightest neutralinos and the lightest chargino, which are almost degenerate in mass even after taking into account radiative corrections. Unfortunately, the collider signals may be very difficult for identifying such a scenario because the pions or leptons in the final state are too soft.

Keywords: Supersymmetry Phenomenology, Hadronic Colliders, Supersymmetric Standard Model.

Contents

1. Introduction

Supersymmetry (SUSY) is one of the leading candidates for physics beyond the standard model (SM). Many of its virtues, including solving the gauge hierarchy problem, gauge coupling unification, dynamical electroweak symmetry breaking, and providing a dark matter (DM) candidate, make it one of the most studied theories. Nevertheless, generic supersymmetric models suffer from various problems such as unsuppressed flavor-changing neutral currents (FCNC), many CP-violating phases, which potentially give rise to large electric dipole moments (EDM's) as well as other CP-violating phenomena, and too many soft SUSY breaking parameters. In order to satisfy the experimental constraints on FCNC and CP violation, the masses of the scalar fermions are pushed to multi-TeV [1], or the CP phases are set at extremely small values or fine-tuned to cancel each other [2]. In particular, if one pushes the first possibility even further so that the scalar fermion masses are extremely large, CP violation and FCNC problems no longer exist. Of course, this re-introduces the fine-tuning problem, and thus one of the motivations for SUSY is lost.

From the landscape point of view in string theories with a huge number of vacua, however, it is not impossible and even more likely to find a vacuum with a high SUSY breaking scale. Based on this observation, Arkani-Hamed and Dimopoulos adopted a rather radical approach to SUSY breaking [3, 4], which was later coined as split SUSY [5]. They essentially discarded the hierarchy problem by accepting the fine-tuning solution to the Higgs boson mass. All the scalars, except for a CP-even Higgs boson, are very heavy. A common mass scale is usually assumed at $\tilde{m} \sim 10^9$ GeV to $M_{\rm GUT}$. However, the gaugino masses M_i and the Higgsino mass parameter μ are comparatively light at the TeV scale, in order to provide an acceptable dark matter candidate [6] and to ensure gauge coupling unification.

Previously, two of us proposed a further splitting in split SUSY by raising the μ parameter to a large value which could be about the same as the sfermion mass or the SUSY breaking scale [7]. It was called the high-μ split SUSY scenario. In such a scenario, we do not encounter the notorious μ problem [8], a viable dark matter candidate is still available, and the gauge coupling unification is only slightly worsened.

In this work, we study a complementary scenario in which only the Higgsino mass parameter μ remains at the weak scale while all other soft SUSY breaking parameters are pushed to $\tilde{m}$ (10^9 GeV). In the framework of minimal supersymmetric standard model (MSSM), this is a minimal scenario since only the SM particle masses and the Higgsino mass parameter μ remain below the weak or TeV scale. The spectrum is very simple: the SM particles, a light CP-even Higgs boson, two neutral Higgsinos and a pair of charged Higgsinos. All the other SUSY particles are at the very high SUSY breaking scale. It can be called the low-μ split SUSY. Although it may be difficult to generate such a scenario from a sensible SUSY breaking model,[1] its phenomenology is so special and unique at colliders that the model deserves a good study, which is the primary goal of the paper. Note that the large separation of scales between the μ parameter and the gaugino masses will induce a radiative correction to the μ parameter such that the μ parameter is within a loop factor of the gaugino mass. We therefore need another fine-tuning between the correction and the bare parameter so that the physical μ parameter is at the electroweak scale while the gaugino masses stay at a high scale.

We summarize the differences between low-μ split SUSY and ordinary split SUSY as follows:

1. The gaugino masses are raised to a very high scale in this scenario while in ordinary split SUSY they are kept at the electroweak scale.

2. The lightest supersymmetric particle (LSP) is the neutral Higgsino, contrary to the ordinary split SUSY model where the Bino, wino, and Higgsino are all possible to be the LSP as the dark matter candidate [11 − 13]

3. In the low-μ split SUSY scenario, the Higgsino dark matter has a negligible elastic scattering cross section with the nuclei in direct detection methods. Instead the Higgsino pair annihilation into gauge boson pairs, particularly a pair of monochromatic photons, is strong. The Bino dark matter in ordinary split SUSY has only a small pair annihilation rate into gauge bosons and diphotons.

[1] Although some models [9, 10] are constructed to give large hierarchies among soft parameters, it is hard to give more than a few orders of magnitude difference.

4. In low-μ split SUSY, the two lightest neutralinos and the lightest chargino have almost degenerate masses. Even the radiative corrections cannot lead to a mass difference more than one GeV between the lightest chargino and the lightest neutralino. Therefore, the decay of the lightest chargino only produces very soft pions or leptons in the final state, which may be too difficult to detect at high energy colliders. In split SUSY, the chargino and neutralino can give rise to interesting signals at hadron and e^+e^- colliders [14 – 16]

Note that the present low-μ split SUSY scenario is rather similar to the focus-point supersymmetry in the phenomenological aspects [17]. Recently, there have been works exploiting the idea of minimal extensions of the SM to satisfy the dark matter and other constraints [18, 19].

The organization of the paper is as follows. In the next section, we examine the gauge coupling unification. In section 3, we describe the mass spectrum of the neutralinos and charginos. Section 4 deals with the couplings relevant to the studies of dark matter and collider phenomenology. In section 5, we discuss the dark matter relic density and the direct and indirect detection. In section 6, we study the phenomenology at hadron and e^+e^- colliders. We conclude in section 7.

2. Gauge coupling unification

The general form of the one-loop renormalization group equations for the gauge couplings between any two mass scales M_X and M_Y is given by

$$\frac{1}{\alpha_i(M_X^2)} = \frac{1}{\alpha_i(M_Y^2)} - \frac{\beta_i}{4\pi} \ln\left(\frac{M_X^2}{M_Y^2}\right) \ , \tag{2.1}$$

where $i = 1, 2, 3$ are indices representing the $U(1)_Y$, $SU(2)_L$, and $SU(3)_C$ gauge couplings, respectively. The differences among the SM, MSSM, ordinary split SUSY, and low-μ split SUSY scenarios reside in the following values of the beta functions :

$$\mathrm{SM} : (\beta)_{\mathrm{SM}} = \begin{pmatrix} 0 \\ -\frac{22}{3} \\ -11 \end{pmatrix} + \begin{pmatrix} \frac{4}{3} \\ \frac{4}{3} \\ \frac{4}{3} \end{pmatrix} F + \begin{pmatrix} \frac{1}{10} \\ \frac{1}{6} \\ 0 \end{pmatrix} N_H \ ,$$

$$\mathrm{MSSM} : (\beta)_{\mathrm{MSSM}} = \begin{pmatrix} 0 \\ -6 \\ -9 \end{pmatrix} + \begin{pmatrix} 2 \\ 2 \\ 2 \end{pmatrix} F + \begin{pmatrix} \frac{3}{10} \\ \frac{1}{2} \\ 0 \end{pmatrix} N_H \ ,$$

$$\text{Split-SUSY} : (\beta)_{\mathrm{split}}|_{<\tilde{m}} = \begin{pmatrix} 0 \\ -6 \\ -9 \end{pmatrix} + \begin{pmatrix} \frac{4}{3} \\ \frac{4}{3} \\ \frac{4}{3} \end{pmatrix} F + \begin{pmatrix} \frac{5}{10} \\ \frac{5}{6} \\ 0 \end{pmatrix} \ ,$$

$$\text{low-}\mu \text{ split SUSY} : (\beta)_{\mu-\mathrm{split}}|_{<\tilde{m}} = \begin{pmatrix} 0 \\ -22/3 \\ -11 \end{pmatrix} + \begin{pmatrix} \frac{4}{3} \\ \frac{4}{3} \\ \frac{4}{3} \end{pmatrix} F + \begin{pmatrix} \frac{5}{10} \\ \frac{5}{6} \\ 0 \end{pmatrix} \ ,$$

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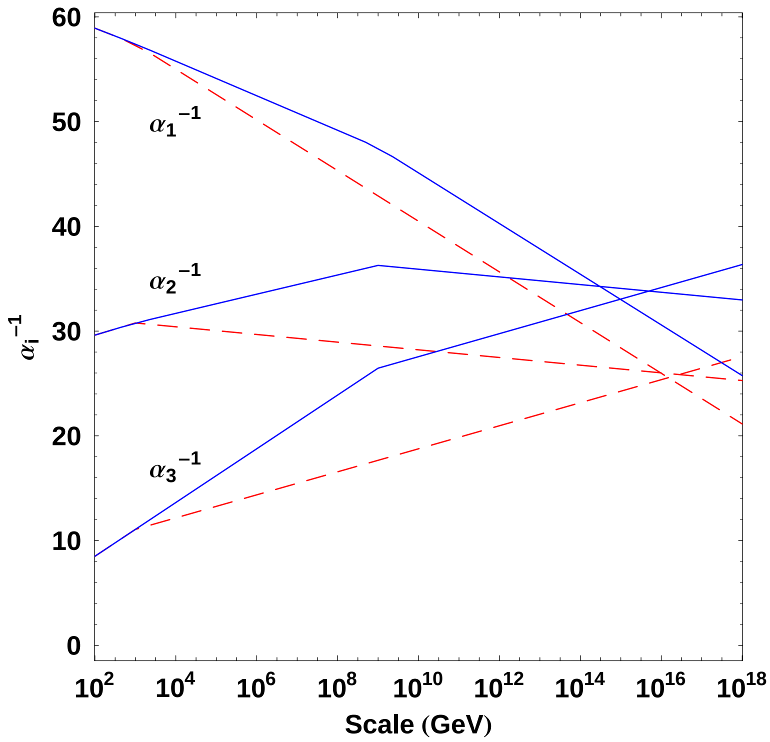

Figure 1: Gauge coupling unification at the one-loop level. In the low-μ split SUSY scenario (indicated by solid curves), the Higgsino masses are set at 1 TeV while all the other soft SUSY parameters at 10^9 GeV. The gauge coupling running in the MSSM (dashed curves) is also included for comparison. Following ref. [3], we take $\alpha_1^{-1}(M_Z) = 58.98$, $\alpha_2^{-1}(M_Z) = 29.57$, and $\alpha_3^{-1}(M_Z) = 8.40$.

where $F = 3$ is the number of generations of fermions or sfermions, and N_H is the number of Higgs doublets ($N_H = 1$ in the SM, $N_H = 2$ in the SUSY). In the evolution of the gauge couplings in the low-μ split SUSY scenario, we use (i) the SM β_i's from the weak scale (m_Z) to μ, the scale of Higgsino masses, which we take a common value of 1 TeV; (ii) the β_i's in low-μ split SUSY scenario from μ to $\tilde{m}$, which we fix at 10^9 GeV; and (iii) the MSSM β_i's from $\tilde{m}$ to the grand unified scale.

The unification of gauge couplings in the low-μ split SUSY case is slightly worse than that in the MSSM, but significantly better than that in the SM. The unification scale is about $10^{14.5}$ GeV. Such a low-scale unification can be dangerous in standard GUTs if there is no other additional symmetries or mechanisms to protect the proton from decaying. One solution is to have a higher dimensional orbifold GUT with a reduced gauge symmetry on the boundary [20]. This also solves the doublet-triplet splitting problem in the usual grand unified models. In this case, threshold effects due to Kaluza-Klein modes have to be included in the running between the compactification scale and the grand unified scale.

If we take the intersection of the α_1 and α_2 curves as the value of gauge coupling strength $\alpha_{\rm GUT}$ at the unification scale (here $\alpha_{\rm GUT} \simeq 1/34.3$, smaller than that in the MSSM) and run it down to the m_Z scale for the strong coupling, we obtain $\alpha_3(m_Z) = 0.098$.

For comparison, the measured value is $\alpha_3(m_Z)_{\rm exp} = 0.1182 \pm 0.0027$ [21] and the one-loop prediction in the MSSM is $\alpha_3(m_Z) = 0.110$. It is well-known that the two-loop effects in the MSSM increase $\alpha_3(m_Z)$ to 0.130. Therefore, we expect that the two-loop contribution in our scenario will lift the predicted $\alpha_3(m_Z)$ closer to the observed value.

To quantify the quality of unification we compute the fractional deviation of $\alpha_3^{-1}(\mathrm{GUT})$ from the $\alpha_{\rm GUT}^{-1}$, which we have taken as the intersecting point of α_1^{-1} and α_2^{-1}. The fractional deviation of $\alpha_3^{-1}(\mathrm{GUT})$ is defined as

$$\frac{\Delta\alpha_3^{-1}(\mathrm{GUT})}{\alpha_{\rm GUT}^{-1}} = \frac{\alpha_3^{-1}(\mathrm{GUT}) - \alpha_{\rm GUT}^{-1}}{\alpha_{\rm GUT}^{-1}} .$$

For consistency we used the 1 loop results. The fractional deviations for MSSM, split SUSY, and low-μ split SUSY are -2.6%, -3.7%, and -5.7%, respectively. Therefore, the quality of unification only gets slightly worse in the low-μ scenario.

3. Mass spectrum

In the low-μ split SUSY scenario, the approximation of $\mu \ll M_1 < M_2$ is very effective. Then the tree-level neutralino masses can be expanded in terms of $\delta = \mu/M_1$:

$$m'_{\widetilde{\chi}_1^0} = \mu\left[1 - \frac{1}{2}(1-s_{2\beta})(c_W^2 + r_2 s_W^2)\frac{r_z^2}{r_2}\delta + \mathcal{O}(\delta^2)\right], \tag{3.1}$$

$$m'_{\widetilde{\chi}_2^0} = -\mu\left[1 + \frac{1}{2}(1+s_{2\beta})(c_W^2 + r_2 s_W^2)\frac{r_z^2}{r_2}\delta + \mathcal{O}(\delta^2)\right], \tag{3.2}$$

$$m'_{\widetilde{\chi}_3^0} = M_1\left[1 + \mathcal{O}(\delta^2)\right], \tag{3.3}$$

$$m'_{\widetilde{\chi}_4^0} = M_2\left[1 + \mathcal{O}(\delta^2)\right], \tag{3.4}$$

where $s_W = \sin\theta_{\rm w}$, $c_W = \cos\theta_{\rm w}$, $s_{2\beta} = \sin 2\beta$, $r_2 = M_2/M_1$, and $r_z = m_Z/\mu$. $\theta_{\rm w}$ is the Weinberg angle, and $\tan\beta$ is the ratio of the two vacuum expectation values in the Higgs sector. The diagonalizing matrix N [22] in the leading order becomes

$$N = \begin{pmatrix} 0 & 0 & \frac{1}{\sqrt{2}} & \mp\frac{1}{\sqrt{2}} \\ 0 & 0 & \pm\frac{1}{\sqrt{2}} & \frac{1}{\sqrt{2}} \\ 1 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \end{pmatrix}, \qquad \text{for } \mu \gtrless 0. \tag{3.5}$$

The chargino masses are

$$m_{\widetilde{\chi}_1^\pm} = \left|\mu\left(1 + s_{2\beta}\frac{m_W^2}{M_2\mu}\right)\right|, \quad m_{\widetilde{\chi}_2^\pm} = M_2\left(1 + s_{2\beta}\frac{m_W^2}{M_2^2}\right), \tag{3.6}$$

which are obtained by the diagonalizing matrices U and V [22] given as, in the leading order,

$$U = \begin{pmatrix} 0 & \pm 1 \\ 1 & 0 \end{pmatrix}, \qquad V = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix} \qquad \text{for } \mu \gtrless 0\,. \tag{3.7}$$

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At tree level, the mass differences with respect to the lightest neutralino are

$$\Delta M_{12}^{\rm tree} \equiv m_{\widetilde{\chi}_2^0} - m_{\widetilde{\chi}_1^0} = \left[1 + t_W^2 \frac{M_2}{M_1}\right] \frac{m_W^2}{M_2}, \tag{3.8}$$

$$\Delta M_{+}^{\rm tree} \equiv m_{\widetilde{\chi}_1^+} - m_{\widetilde{\chi}_1^0} = \frac{1}{2}\left[1 + (1 - s_{2\beta})\, t_W^2 \frac{M_2}{M_1}\right] \frac{m_W^2}{M_2}\,, \tag{3.9}$$

where $t_W = s_W/c_W$. It is easy to see that if $\tan\beta \gg 1$, $\Delta M_{12}^{\rm tree} \simeq 2\,\Delta M_{+}^{\rm tree}$. For $M_1 \sim M_2 \sim 10^8 - 10^9$ GeV, $\Delta M_{12}^{\rm tree} \simeq 2\,\Delta M_{+}^{\rm tree} \sim \mathcal{O}(10 - 100)$ keV.

With heavy gaugino masses at the order of 10^9 GeV, the mass differences due to radiative corrections are more important than the tree-level mass differences. The one-loop radiative mass corrections to neutralinos and charginos were performed in ref. [23], where the complete one-loop self-energies for charginos and neutralinos are included in their mass matrices. The total mass matrices are diagonalized to obtain the final mass spectrum. The one-loop neutralino mass matrix is

$$\mathcal{M}_N = \mathcal{M}_N^{(0)} + \frac{1}{2}\left[\delta\mathcal{M}_N(p^2) + \delta\mathcal{M}_N^T(p^2)\right]\,, \tag{3.10}$$

and the one-loop chargino mass matrix is

$$\mathcal{M}_C = \mathcal{M}_C^{(0)} - \Sigma_R^+(p^2)\mathcal{M}_C - \mathcal{M}_C\Sigma_L^+(p^2) - \Sigma_S^+(p^2). \tag{3.11}$$

The detailed expressions for $\mathcal{M}_N$ and $\mathcal{M}_C$ are referred to ref. [23]. We note that the one-loop neutralino mass matrix is still symmetric.

The mass degeneracy between $\widetilde{\chi}_1^\pm$ and $\widetilde{\chi}_1^0$ is lifted by the radiative corrections, given by

$$\Delta M_+^{\rm rad} = \frac{g^2}{8\pi^2} s_W^2 \mu \left[B_1(\mu,\mu,m_Z) - 2B_0(\mu,\mu,m_Z) - \{B_1(\mu,\mu,0) - 2B_0(\mu,\mu,0)\}\right]\,, \tag{3.12}$$

where we have neglected the contributions from super heavy particles like sfermions and non-SM Higgs bosons. The functions B_0 and B_1 are defined in [23, appendix B].

Using the relation

$$2\big[B_1(M,M,m) - 2B_0(M,M,m)\big] = \frac{1}{\epsilon} + f\left(\frac{m}{M}\right), \tag{3.13}$$

the mass difference is simplified to

$$\Delta M_+^{\rm rad} = \frac{\alpha_2 \mu}{4\pi} s_W^2 \left[f\left(\frac{m_Z}{\mu}\right) - f(0)\right]. \tag{3.14}$$

The function $f(a)$ is [24]

$$f(a) = \int_0^1 dx\, 2(1+x)\ln(x^2 + (1-x)a^2) \simeq -5 + c\,a + \mathcal{O}(a^2), \tag{3.15}$$

in which the second equality holds for $a \ll 1$. Numerically, c is about 6.3. For $\mu \gg m_Z$, we have $\Delta M_+^{\rm rad}$ almost independent of μ, given by

$$\Delta M_+^{\rm rad} = \frac{\alpha_2}{4\pi} c\, m_Z\, s_W^2 \quad \text{for } \mu \gg m_Z\,. \tag{3.16}$$

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Numerically it is about 340 MeV, much larger than the tree-level splitting for $M_{1,2} \simeq 10^9$ GeV. For comparison, we present the radiative mass difference in the wino-LSP scenario [24]:

$$\Delta M_+^{\rm rad}|_{\rm wino} = \frac{\alpha_2}{4\pi} c\, m_Z\, c_W(1-c_W) \quad \text{for } \mu \gg m_Z. \tag{3.17}$$

Since $c_W(1-c_W) \simeq 0.11$ for $s_W^2 \simeq 0.23$, the radiative mass correction for the Higgsino case is about twice as much as the wino case. Therefore, the lightest chargino will decay into the lightest neutralino plus a soft pion, or a charged lepton with a neutrino. We cannot avoid the decay of the lightest chargino into a pion.

On the other hand, $\Delta M_{12}(\equiv m_{\widetilde{\chi}_2^0} - m_{\widetilde{\chi}_1^0})$ remains intact from the radiative mass corrections in the low-μ split SUSY case. The radiative corrections modify the neutralino mass matrix in the way preserving the symmetry property, $\mathcal{M}_N = \mathcal{M}_N^T$. Neglecting the extremely small μ/M_1 correction, $\Delta M_{12}^{\rm rad}$ comes from the difference between the $(\mathcal{M}_N)_{34}$ and $(\mathcal{M}_N)_{43}$ components, which are the same. $\widetilde{\chi}_1^0$ and $\widetilde{\chi}_2^0$ are, therefore, highly degenerate with a mass difference of the order of $(10-100)$ keV in the low-μ split SUSY scenario. This small difference prohibits the decay channel of $\widetilde{\chi}_2^0 \to \widetilde{\chi}_1^0 Z^* \to \widetilde{\chi}_1^0 e^+e^-$. The only open modes are $\widetilde{\chi}_2^0 \to \widetilde{\chi}_1^0 Z^* \to \widetilde{\chi}_1^0 \nu\bar{\nu}$ and $\widetilde{\chi}_2^0 \to \widetilde{\chi}_1^0 \gamma$. Unfortunately both are difficult to probe, as the former will be purely the missing energy signal and the latter will generate too soft photons with missing energy.

4. Relevant couplings

As the gaugino mass parameters M_1, M_2 and M_3 become very large while the μ parameter remains light, only the two lightest neutralinos and the lightest chargino are accessible at high energy colliders. In this section, we highlight the couplings that are relevant to studies on collider and dark matter phenomenology. Let us first examine their relevant couplings to gauge bosons and the Higgs bosons.

- The Z-$\widetilde{\chi}_{1,2}^0$-$\widetilde{\chi}_{1,2}^0$ couplings only receive contributions from the Higgsino-Higgsino-gauge couplings. In the limit of very large M_1 and M_2 the Higgsino component of $\widetilde{\chi}_1^0$ and $\widetilde{\chi}_2^0$ are essentially one. Therefore, the Z boson couplings to the neutralinos are large.

- The H-$\widetilde{\chi}_{1,2}^0$-$\widetilde{\chi}_{1,2}^0$ couplings have sources from the Higgs-Higgsino-gaugino terms. Therefore, in the limit of large M_1 and M_2, these couplings go to zero.

- The W^--$\widetilde{\chi}_{1,2}^0$-$\widetilde{\chi}_1^+$ couplings have sources from the Higgsino-Higgsino-gauge couplings and from the gaugino-gaugino-gauge couplings. In the limit of large M_1 and M_2, the latter contribution goes to zero while the former remains. Therefore, the W^--$\widetilde{\chi}_{1,2}^0$-$\widetilde{\chi}_1^+$ couplings contain only the Higgsino-Higgsino-gauge part.

- The H^--$\widetilde{\chi}_{1,2}^0$-$\widetilde{\chi}_1^+$ couplings have sources from the Higgs-Higgsino-gaugino couplings. In the limit of large M_1 and M_2, they do not contribute to H^--$\widetilde{\chi}_{1,2}^0$-$\widetilde{\chi}_1^+$, which thus vanish.

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- The $\gamma(Z)$-$\widetilde{\chi}_1^+$-$\widetilde{\chi}_1^-$ couplings have sources from the Higgsino-Higgsino-gauge couplings and from the gaugino-gaugino-gauge couplings. In the limit of large M_1 and M_2, the latter contribution goes to zero while the former remains. Therefore, the $\gamma(Z)$-$\widetilde{\chi}_1^+$-$\widetilde{\chi}_1^-$ coupling contains only the Higgsino-Higgsino-gauge part.

- The H-$\widetilde{\chi}_1^+$-$\widetilde{\chi}_1^-$ couplings have sources from the Higgs-Higgsino-gaugino couplings. In the limit of large M_1 and M_2, they do not contribute to H-$\widetilde{\chi}_1^+$-$\widetilde{\chi}_1^-$, which thus vanish.

The only couplings that survive in the low-μ split SUSY scenario are Z-$\widetilde{\chi}_{1,2}^0$-$\widetilde{\chi}_{1,2}^0$, W^--$\widetilde{\chi}_{1,2}^0$-$\widetilde{\chi}_1^+$, and $\gamma(Z)$-$\widetilde{\chi}_1^+$-$\widetilde{\chi}_1^-$. The phenomenology of the two light neutralinos and the lightest chargino depends on the interaction Lagrangian

$$\begin{aligned}\mathcal{L} = &\frac{g}{2\cos\theta_{\rm w}}\overline{\widetilde{\chi}_i^0}\gamma^\mu\left(O_{ij}^{L''}P_L + O_{ij}^{R''}P_R\right)\widetilde{\chi}_j^0\, Z_\mu \\ &+\left[g\overline{\widetilde{\chi}_i^0}\gamma^\mu\left(O_{ij}^{L}P_L + O_{ij}^{R}P_R\right)\widetilde{\chi}_j^+\, W_\mu^- + H.c.\right] \\ &- e\overline{\widetilde{\chi}_i^+}\gamma^\mu\widetilde{\chi}_i^+\, A_\mu - \frac{g}{\cos\theta_{\rm w}}\overline{\widetilde{\chi}_i^+}\gamma^\mu\left(O_{ij}^{L'}P_L + O_{ij}^{R'}P_R\right)\widetilde{\chi}_j^+\, Z_\mu\,, \end{aligned} \tag{4.1}$$

where

$$\begin{aligned} O_{ij}^{L''} &= \frac{1}{2}\left(N_{i4}N_{j4}^* - N_{i3}N_{j3}^*\right), \qquad O_{ij}^{R''} = -\left(O_{ij}^{L''}\right)^*, \\ O_{ij}^{L} &= N_{i2}V_{j1}^* - \frac{1}{\sqrt{2}}N_{i4}V_{j2}^*, \qquad O_{ij}^{R} = N_{i2}^*U_{j1} + \frac{1}{\sqrt{2}}N_{i4}^*U_{j2}, \\ O_{ij}^{L'} &= \cos^2\theta_{\rm w}V_{i1}V_{j1}^* + \frac{\cos 2\theta_{\rm w}}{2}V_{i2}V_{j2}^*, \qquad O_{ij}^{R'} = \cos^2\theta_{\rm w}U_{j1}U_{i1}^* + \frac{\cos 2\theta_{\rm w}}{2}U_{j2}U_{i2}^*. \end{aligned}$$

The mixing matrices in eqs. (3.5) and (3.7) further simplify the couplings as

$$\begin{aligned} \text{for } \mu \gtrless 0,\quad & O_{11}^{L''} = O_{22}^{L''} = 0, \quad O_{12}^{L''} = \mp\frac{1}{2}, \\ & O_{11}^{L} = \pm\frac{1}{2}, \quad O_{11}^{R} = -\frac{1}{2}, \\ & O_{11}^{L'} = O_{11}^{R'} = \frac{1}{2}\cos 2\theta_{\rm w}. \end{aligned} \tag{4.2}$$

5. Dark Matter

5.1 Relic Density

The Higgsino LSP is well-known to have large annihilation cross sections into ZZ and WW pairs and also into $f\bar{f}$ via Z boson exchange. In addition, the lightest neutralino is close in mass with the second lightest neutralino and the lightest chargino. Such large annihilation cross sections and effective co-annihilation reduce significantly the thermal relic density of the Higgsino LSP. The simplified formula for the Higgsino LSP is given by

$$\Omega_{\tilde{H}}h^2 \simeq 0.1\left(\frac{M_{\rm LSP}}{1\,{\rm TeV}}\right)^2.$$

Only heavy neutral Higgsino with mass ~ 1 TeV can explain the dominant dark matter if thermal production is the only source.[2] Unfortunately, in this case SUSY particles can only be marginally produced at the LHC.

On the other hand, there can be non-thermal sources of dark matter. In such a case, the mass of the Higgsino can be lower. Both collider and dark matter experiments can find interesting signals. A good example of non-thermal sources can be found in anomaly-mediated SUSY breaking models where the neutral wino is the LSP [26, 27]. For a relatively light neutral wino it cannot be the dominant dark matter because of its large annihilation cross sections. However, an intriguing source of non-thermal wino for compensation is the decay of moduli fields [28], which can produce a sufficient amount of neutral winos. This case is similar to the Higgsino dark matter. There are also other non-thermal sources of Higgsino or wino dark matter discussed in the literatures [29], such as the gravitino NLSP decay and decay of some scalar relics.

These non-thermal sources have important constraints from cosmology; their late decay should be before the big bang nucleosynthesis (BBN) and the reheating temperature due to the decay of the moduli should be above MeV, so as not to disturb the success of the BBN. The moduli decay when the expansion rate of the Universe becomes compatible with the decay width of the moduli. In the model used in ref. [28], the reheating temperature being above MeV requires the mass of the modulus field to be above 100 TeV. Therefore, as long as the decay of moduli is before BBN and the reheating temperature is above MeV, it is consistent with BBN.

We also note that dark matter can be made up of some almost non-interacting particles, e.g., axions, which have nothing to do with the electroweak scale physics. An interesting scenario of this kind [30] is proposed, in which all sparticles are super-heavy and the dark matter is explained by the axion.

5.2 Direct dark matter detection

In general, the neutralino LSP interacts with the nucleon of the detecting material via Higgs boson exchange and/or squark exchange for the spin-independent cross section while via Z boson exchange and part of the squark exchange for spin-dependent cross section. Since the signal is very mild, a large coherence effect in the detecting material is crucial, which is only useful for the spin-independent cross section. In the low-μ split SUSY scenario, however, the Higgs-Higgsino-gaugino couplings are zero because of the decoupling of the gauginos, and the squarks are extremely heavy. Thus, the spin-independent cross section is negligible [31].

Here we concentrate on the spin-dependent cross sections. Even though the current detector sensitivity is much lower, there are some proposals that focus on measuring the spin-dependent cross sections [32]. In split SUSY, the contributions from super-heavy squarks are negligible. Therefore, the only remaining contribution comes from the Z-boson exchange via the Higgsino-Higgsino-gauge type coupling.

[2]There was a light Higgsino window, where $m_{\widetilde{\chi}_1^0} < m_W$, such that the annihilation $\widetilde{\chi}_1^0\widetilde{\chi}_1^0 \to W^+W^-$ is suppressed and can satisfy the relic density constraint. However, the present limit on chargino mass from LEP2 [25] rules out this window. The limit is $m_{\widetilde{\chi}_1^+} > 103$ GeV for $\Delta m \equiv m_{\widetilde{\chi}_1^+} - m_{\widetilde{\chi}_1^0} < 100$ MeV or > 3 GeV, while the limit is weaken to $m_{\widetilde{\chi}_1^+} > 92$ GeV for $\Delta m \sim 0.1 - 3$ GeV.

Let us remind the reader of the spin-dependent cross section with protons and neutrons:

$$\sigma_{\chi p}^{SD} = \frac{3\mu_{\chi p}^2}{\pi} |G_a^p|^2 , \qquad \sigma_{\chi n}^{SD} = \frac{3\mu_{\chi n}^2}{\pi} |G_a^n|^2 , \tag{5.1}$$

where $\mu_{\chi p}$ ($\mu_{\chi n}$) is the reduced mass of the neutralino and proton (neutron), and the effective axial couplings $G_a^{p,n}$ are

$$G_a^{p,n} = \sum_{u,d,s} (\Delta q)_{p,n} \frac{g_{Z\chi\chi} g_{Zqq}}{m_Z^2} , \tag{5.2}$$

where

$$g_{Z\chi\chi} = \frac{g}{2\cos\theta_{\rm w}} \left(|N_{13}|^2 - |N_{14}|^2\right) , \qquad g_{Zqq} = -\frac{g}{2\cos\theta_{\rm w}} T_{3q} .$$

It is easy to see that in this low-μ split SUSY scenario ($\mu \ll M_1, M_2$) the H_u^0 and H_d^0 mix maximally so that $|N_{13}| = |N_{14}| = 1/\sqrt{2}$. Therefore, the vanishing factor $N_{13}^2 - N_{14}^2 \simeq 0$ suppresses the spin-dependent cross section also.

Since the mass difference between $\widetilde{\chi}_1^0$ and $\widetilde{\chi}_2^0$ is of the order of $10 - 100$ keV for $M_2 \sim 10^9 - 10^8$ GeV, there is a slight chance that the $\widetilde{\chi}_1^0$ transits into $\widetilde{\chi}_2^0$ after scattering with the nuclei in the detector, via the nonzero Z-$\widetilde{\chi}_1^0$-$\widetilde{\chi}_2^0$ coupling [33]. The velocity of the weakly interacting massive particle (WIMP) in our galaxy follows a Boltzmann distribution centered at $v = 270$ km s$^{-1} \sim 10^{-3}c$. The spectrum of recoil is exponential with a typical energy $\langle E \rangle \sim 50$ keV. Therefore, if $M_2 > 10^9$ GeV and so the mass difference $\Delta M_{12} \lesssim$ a few keV, $\widetilde{\chi}_1^0$ can transit into $\widetilde{\chi}_2^0$ after the scattering. This is because the recoil energy is large enough for the transition to take place. While this could be an interesting signal [33], we require $M_2 \lesssim 10^9$ GeV so that such transitions cannot take place. Moreover, the gluino cosmology requires a SUSY breaking scale to be less than 10^9 GeV [34]. Therefore, the direct detection experiments do not have constraints on the present scenario.

Although the tree-level spin-independent and spin-dependent scattering cross sections are negligible in this scenario, one-loop corrections give a non-zero cross section [35]. The cross section is at the 10^{-45} cm^2 level [35], which is about two orders of magnitude below the current best limit of 10^{-43} cm^{-2} of CDMS II [36].

5.3 Indirect dark matter detection

The Higgsino dark matter can have very interesting signals for indirect detection in view of its large annihilation cross sections into ZZ and W^+W^- pairs, as well as into $\gamma\gamma, \gamma Z$ via one-loop diagrams. In particular, the last two channels, though loop-suppressed, can give a very clean signal of monochromatic photon lines. If the resolution of the photon detectors (either ground-based or satellite-based) is high enough, a clean and unambiguous photon peak at hundreds of GeV can be observed above the background.

Here we give an estimate of the photon flux in the low-μ split SUSY scenario in which only the W^--$\widetilde{\chi}_1^+$ loop is important. Using the results given in ref. [37], we obtain

$$v\sigma(\widetilde{\chi}_1^0\widetilde{\chi}_1^0 \to \gamma\gamma) \simeq 1 \times 10^{-28}\ {\rm cm^3 s^{-1}} , \tag{5.3}$$

which is roughly independent of the mass of the neutralino from a few hundred GeV to a few TeV. For comparison, the value of $v\sigma$ for a pure wino case is about 14×10^{-28} cm^{3}s^{-1} [7]. The photon flux as a result of this annihilation is given by [38]

$$\begin{aligned}\Phi_\gamma &\simeq 1.87 \times 10^{-11} \left(\frac{N_\gamma v\sigma}{10^{-29}\,\mathrm{cm}^3\mathrm{s}^{-1}}\right)\left(\frac{10\,\mathrm{GeV}}{M_{\widetilde{\chi}_1^0}}\right)^2 J(\psi)\ \mathrm{cm}^{-2}\mathrm{s}^{-1}\mathrm{sr}^{-1} \\ &\simeq 1.5 \times 10^{-11}\ \mathrm{cm}^{-2}\mathrm{s}^{-1}\mathrm{sr}^{-1}\,, \end{aligned} \tag{5.4}$$

where we have used $v\sigma = 1 \times 10^{-28}\,\mathrm{cm}^3\mathrm{s}^{-1}$, $M_{\widetilde{\chi}_1^0} = 500$ GeV, $N_\gamma = 2$, and $J(\psi = 0) = 100$ for the photon flux coming from the Galactic Center. The value of $J(\psi)$ depends on the selected Galactic halo model, which ranges from $O(10)$ to $O(1000)$ [38]. For a typical Atmospheric Cerenkov Telescope (ACT) such as VERITAS [39] and HESS [40], the angular coverage is about $\Delta\Omega = 10^{-3}$ and may reach a sensitivity at the level of $(10^{-14} - 10^{-13})\,\mathrm{cm}^{-2}\,\mathrm{s}^{-1}$ at the TeV scale. Therefore, the signal of pure Higgsino dark matter annihilating into monochromatic photons is easily covered by the next generation ACT experiments. Note that there is a nonperturbative enhancement to the annihilation due to the bound-state effect by exchanging the $SU(2)_L$ partners [41]. The enhancement is significant when the LSP is heavier than one TeV and is nearly degenerate with its $SU(2)_L$ partners, as in the scenarios of wino- and Higgsino-LSP. However, in our case of degenerate Higgsino-LSP (doublet case) the enhancement at around 1 TeV is only within a factor of two [41], which is less than the uncertainty in the dark matter density profile.

Since the Higgsino annihilation into the W^+W^- and ZZ pairs is very effective, one can also measure the excess in anti-protons and positrons [42], which can be measured in anti-matter search experiments, e.g., AMSII [43], as well as the neutrino flux from the core of the Sun.

6. Collider phenomenology

The collider phenomenology in the low-μ split SUSY scenario is mainly concerned with the production and detection of neutralinos and charginos. We restrict our discussions below to the case with exact or approximate R-parity symmetry.

6.1 Gaugino pair production

Note that the masses of the two lightest neutralinos and the lightest chargino come from a single parameter μ. They are almost degenerate in mass at tree level, and the mass splitting is of $O(10)$ keV, as shown in section 3. The radiative corrections lift the mass degeneracy such that the lightest chargino is slightly heavier than the lightest neutralino by a mass splitting at least 340 MeV.

The only pair production channels at hadronic colliders are $\widetilde{\chi}_1^0\widetilde{\chi}_2^0$, $\widetilde{\chi}_1^+\widetilde{\chi}_1^-$, and $\widetilde{\chi}_{1,2}^0\widetilde{\chi}_1^\pm$ through the Drell-Yan process. Note that vanishing couplings of Z-$\widetilde{\chi}_1^0$-$\widetilde{\chi}_1^0$ and Z-$\widetilde{\chi}_2^0$-$\widetilde{\chi}_2^0$ suppress the production of $\widetilde{\chi}_1^0\widetilde{\chi}_1^0$ and $\widetilde{\chi}_2^0\widetilde{\chi}_2^0$. Since the only relevant parameter is μ, we expect that their cross sections show fixed ratios. In figure 2 we present the production cross sections versus the μ parameter for the Tevatron and LHC. We sum the cross sections

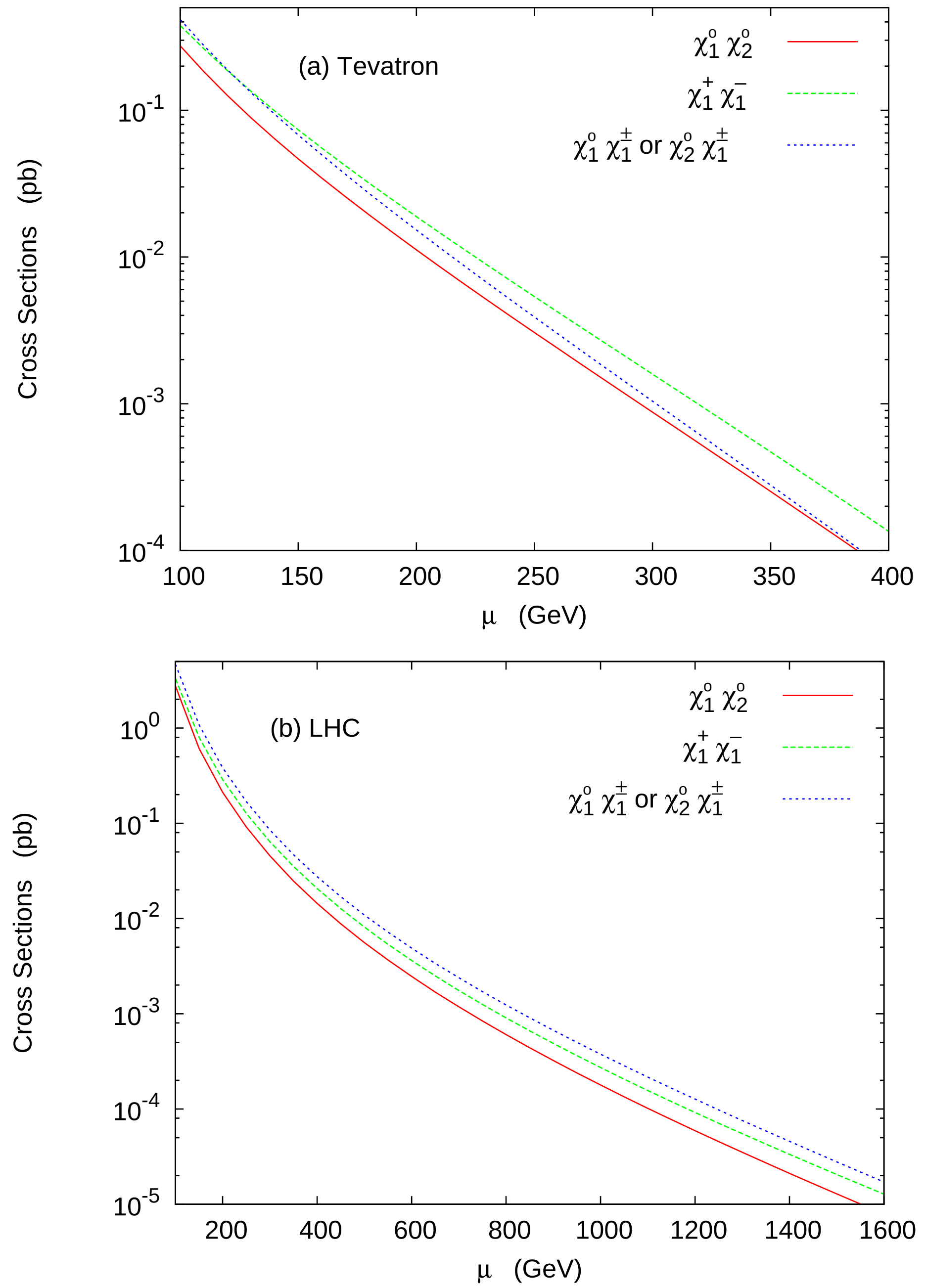

Figure 2: Production cross sections versus μ (the Higgsino mass parameter at the weak scale) for the $\widetilde{\chi}_1^0\widetilde{\chi}_2^0$, $\widetilde{\chi}_1^+\widetilde{\chi}_1^-$, $\widetilde{\chi}_1^0\widetilde{\chi}_1^\pm$ ($\pm$ states summed), and $\widetilde{\chi}_2^0\widetilde{\chi}_1^\pm$ channels at (a) the Tevatron and (b) the LHC.

of $\widetilde{\chi}_1^0\widetilde{\chi}_1^+$ and $\widetilde{\chi}_1^0\widetilde{\chi}_1^-$ in the figure. Note that at the Tevatron ($p\bar{p}$ collision at 1.96 TeV), the production cross sections for $\widetilde{\chi}_1^0\widetilde{\chi}_1^+$ and $\widetilde{\chi}_1^0\widetilde{\chi}_1^-$ are the same. Furthermore, since $\widetilde{\chi}_1^0$ and $\widetilde{\chi}_2^0$ are almost degenerate in mass, the curves for the production cross sections of $\widetilde{\chi}_1^0\widetilde{\chi}_1^\pm$ and $\widetilde{\chi}_2^0\widetilde{\chi}_1^\pm$ are indistinguishable in the figure. At the LHC, the sum of cross sections for $\widetilde{\chi}_1^+\widetilde{\chi}_1^-$ and $\widetilde{\chi}_1^\pm\widetilde{\chi}_{1,2}^0$ is of the order of 1 fb for $\mu = 1$ TeV (the relevant scale of dark matter).

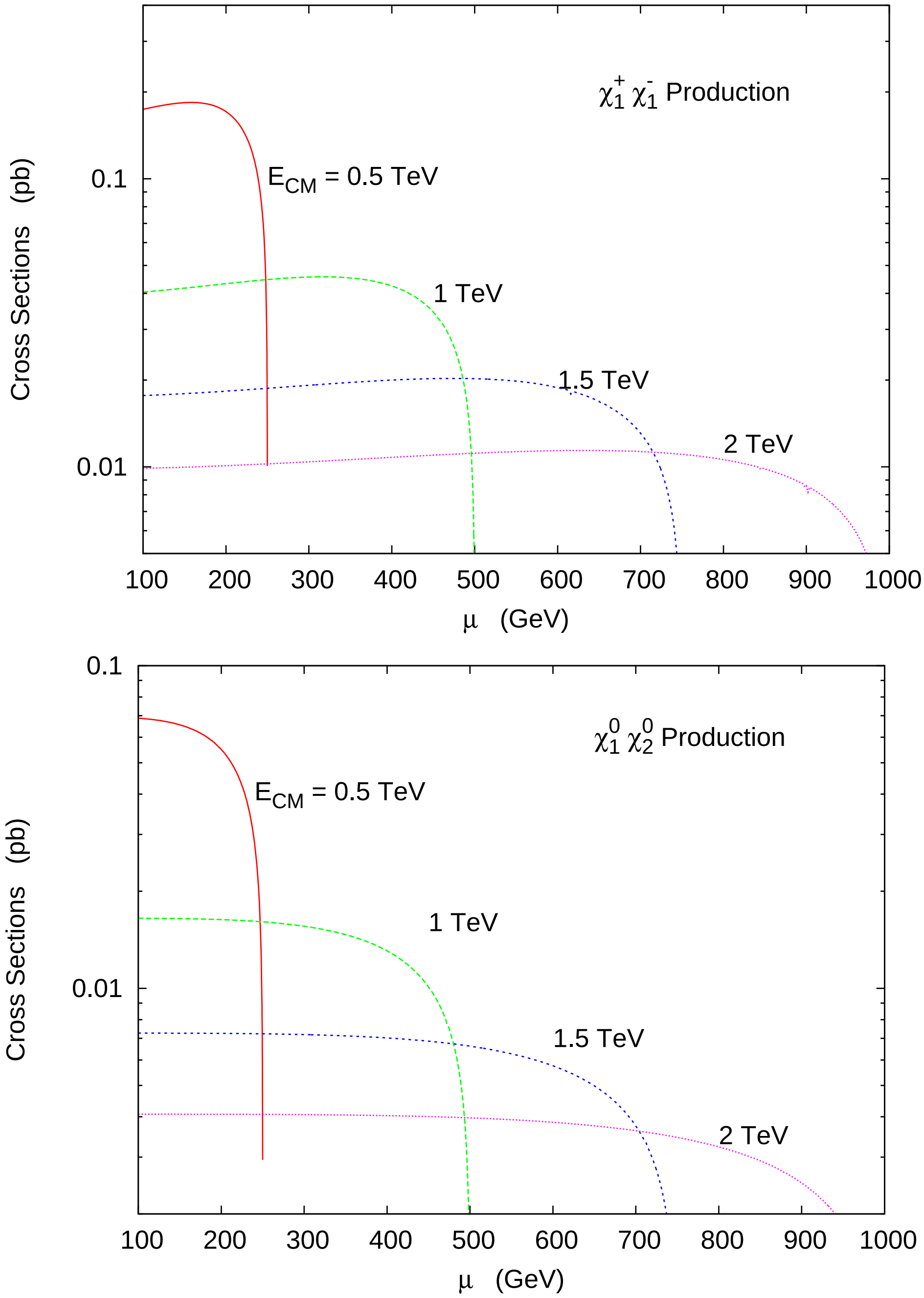

Figure 3: Production cross sections versus μ (the Higgsino mass parameter at the weak scale) for the $\widetilde{\chi}_1^+\widetilde{\chi}_1^-$ and $\widetilde{\chi}_1^0\widetilde{\chi}_2^0$ at e^+e^- linear collider for $\sqrt{s} = 0.5, 1, 1.5, 2$ TeV.

This is a reasonable cross section given an integrated $O(100)$ fb^{-1} luminosity, but the most important criterion is the detection efficiency. The leptons or pions from the chargino decay are too soft for efficient detection. Details will be given in the next subsection. On the other hand, with the clean environment at e^+e^- collider one may be able to readily identify such events using the initial-state radiation [25].

At e^+e^- linear colliders, one can consider neutralino pair $\widetilde{\chi}_1^0\widetilde{\chi}_2^0$ and chargino pair $\widetilde{\chi}_1^+\widetilde{\chi}_1^-$ production. We show their cross sections in figure 3. Since the μ parameter is the only parameter, the cross sections show interesting ratios, given by the gauge coupling.

6.2 Neutralino and chargino decays

In the case of the lightest neutralino as the LSP, the lightest chargino will decay into the lightest neutralino plus leptons or jets. Even though the mass degeneracy is lifted by radiative corrections, the mass difference is quite small and only of the order of twice the pion mass. The decay products are thus very soft. Experimentally, it is a challenge to tag these soft leptons or jets. The decay rate of the chargino into the neutralino and a virtual W boson depends critically on the mass difference $\Delta M_+ \equiv m_{\widetilde{\chi}_1^+} - m_{\widetilde{\chi}_1^0}$. The phenomenology in this case had been studied in great detail in ref. [44]. We give highlights as follows.

The partial decay width of $\widetilde{\chi}_1^+ \to \widetilde{\chi}_1^0 f\bar{f}'$ is given by [44]

$$\begin{aligned}
&\Gamma(\widetilde{\chi}_1^+ \to \widetilde{\chi}_1^0 f\bar{f}') \\
&\quad = \frac{N_c G_F^2}{(2\pi)^3}\Bigg\{M_{\widetilde{\chi}_1^+}\left[\left(O_{11}^L\right)^2 + \left(O_{11}^R\right)^2\right] \\
&\qquad\qquad \times \int_{(M_{\widetilde{\chi}_1^0}+m_f)^2}^{M^2_{\widetilde{\chi}_1^+}} dq^2 \left(1 - \frac{M^2_{\widetilde{\chi}_1^0} + m_f^2}{q^2}\right)\left(1 - \frac{q^2}{M^2_{\widetilde{\chi}_1^+}}\right)^2 \lambda^{1/2}(q^2, M^2_{\widetilde{\chi}_1^0}, m_f^2) \\
&\qquad\qquad -2M_{\widetilde{\chi}_1^0} O_{11}^L O_{11}^R \int_{m_f^2}^{\Delta M_+^2} dq^2\, \frac{q^2}{M^2_{\widetilde{\chi}_1^+}}\left(1 - \frac{m_f^2}{q^2}\right)^2 \lambda^{1/2}(M^2_{\widetilde{\chi}_1^+}, M^2_{\widetilde{\chi}_1^0}, q^2)\Bigg\}\,, \qquad (6.1)
\end{aligned}$$

where (f, f') is, for example, (u, d) or (e, ν_e), $N_c = 3(1)$ if f is a quark (lepton), and $\lambda(a, b, c) = (a+b-c)^2 - 4ab$. In low μ-split SUSY, $|O_{11}^L| = |O_{11}^R| = 1/2$. The above formula is valid for (i) leptonic decays and (ii) hadronic decays when $\Delta M_+ \gtrsim 2$ GeV. For hadronic decays with $\Delta M_+ \lesssim 1 - 2$ GeV, one has to explicitly sum over exclusive hadronic final states. We have to include the partial decay widths for the decays into one, two, and three pions. The explicit formulas can be found in ref. [44].

Generally speaking, the detection of the chargino depends on the size of ΔM_+:

1. $\Delta M_+ < m_\pi$. As we have explained before, this case will not happen due to the one-loop radiative corrections that result in a mass splitting greater than the pion mass.

2. $m_\pi < \Delta M_+ < 1$ GeV. This is the most difficult regime to probe experimentally, and very much depends on the design of the central detector. Important criteria are the decay length $c\tau$ of the chargino and the momentum of the pion from the chargino decay. The decay length $c\tau$ may be long enough for the chargino to travel through a few layers of the silicon vertex detector. For example, if $m_\pi < \Delta M_+ < 190$ MeV the chargino will typically pass through at least two layers of silicon chips [44]. Since the pion is produced from the chargino decay, it is a non-pointing pion; the

backward extrapolation of the pion track does not lead to the interaction point. The resolution on the impact parameter $b_{\rm res}$ depends on the momentum of the pion $p_\pi \sim \sqrt{\Delta M_+^2 - m_\pi^2}$ in the chargino rest frame. The higher the momentum is, the better the resolution $b_{\rm res}$ will be [44]. Thus, detecting the signal involves the combination of detecting a track left in only a few layers of the silicon detectors and identifying a nonzero impact parameter of the pion coming out of the chargino. A detailed simulation is beyond the scope of the present paper.

3. $\Delta M_+ \gtrsim 1 - 2$ GeV. We can use eq. (6.1) to estimate the total decay width of the chargino. The decay width is large enough that the decay is prompt, producing soft leptons, pions, or jets, plus large missing energy. The problem is on the softness of the decay products, whose detection is experimentally difficult. Only when ΔM_+ is sufficiently large to produce hard enough leptons or jets can the chargino decay be detected. Otherwise, one has to rely on some other channels, such as $e^+e^- \to \gamma\widetilde{\chi}_1^+\widetilde{\chi}_1^- \to \gamma + \not{E}_T$, a single photon plus large missing energy above the SM background $e^+e^- \to \gamma\nu\bar{\nu}$ [44]. Unfortunately, the signal rate is $\mathcal{O}(\alpha_{\rm em})$ smaller than the chargino pair production. Detecting such a signal is even more difficult at hadronic colliders.

In summary, the detection of the chargino is easier only when ΔM_+ is larger than a few GeV. As we show in section 3, ΔM_+ from the one-loop corrections is at least 340 MeV. The intermediate range between this value and 1 GeV poses a challenge to experiments. The questions are how many layers of silicon the chargino can travel and how well the resolution of the non-pointing pion can be.

On the other hand, the mass splitting between the second lightest and the lightest neutralino is even smaller than the electron mass when $\tilde{m} \gtrsim 10^{8-9}$ GeV. Only $\widetilde{\chi}_2^0 \to \widetilde{\chi}_1^0\gamma$ and $\widetilde{\chi}_2^0 \to \widetilde{\chi}_1^0\nu\bar{\nu}$ are possible. The former decay mode produces a soft photon with keV energy while the latter leads to a totally missing energy signal. In addition, the decay length is of the order of $10^9 - 10^{12}$ m for $\tilde{m} = 10^8 - 10^9$ GeV. Therefore, it is impossible to detect this second lightest neutralino as long as the mass splitting ΔM_{12} is less than twice the electron mass. We will have an excess of the totally missing energy signal over the SM background of $e^+e^- \to \nu\bar{\nu}$ at e^+e^- colliders. In the case of $\tilde{m} < 10^7$ GeV, the mass splitting ΔM_{12} is larger than the electron mass. The charged lepton mode is open but the leptons are still too soft for identification.

In this almost hopeless scenario, one has to rely on additional tags of the processes, such as $e^+e^- \to \widetilde{\chi}_1^+\widetilde{\chi}_1^-\gamma$ at e^+e^- colliders [47] or $q\bar{q} \to \widetilde{\chi}_1^+\widetilde{\chi}_1^- g\,(\gamma)$, $\widetilde{\chi}_1^+\widetilde{\chi}_{1,2}^0 g\,(\gamma)$ at hadron colliders.

7. Conclusions

In the present paper we have considered a low-μ split SUSY scenario. The only parameter beyond the SM is the Higgsino mass parameter μ, which gives rise to two light neutralinos and a pair of light charginos. We summarize its characteristic features as follows:

1. Gauge coupling unification is as good as that in MSSM or split SUSY. The unification point is lower at 10^{14-15} GeV.

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2. The neutral Higgsino is the LSP. The mass degeneracy among the two lightest neutralinos and the lightest chargino gives a very large coannihilation effect such that the thermal source of Higgsino dark matter dominates only if the Higgsino mass is at least 1 TeV. On the other hand, it can have other non-thermal sources, which involve other unknown parameters of the model.

3. The Higgsino dark matter has negligible direct detection rates, which only arise from loop corrections. However, the pair annihilation cross sections into $WW, ZZ, \gamma\gamma$ are large. Thus, we can look for positron or anti-proton excess in nearby galaxies, as well as monochromatic photon lines from the Galactic Center.

4. The collider phenomenology is also quite different from the usual MSSM or split SUSY. The only extra parameter beyond the SM is the μ parameter. The electroweak gaugino pair production shows a fixed ratio because they are produced via the gauge couplings.

5. The decays of the second lightest neutralino and the lightest chargino are difficult to detect because the chargino decay length is not long enough and its decay products are too soft. One has to rely on additional tags of the processes.

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Note added : during the write-up a couple of works [45, 46] similar to the present one appears. We have different discussions with the former paper. We have small overlaps with the second one.

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