# Split-SUSY dark matter in light of direct detection limits (Postprint)

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## Full Text

## Split-SUSY Dark Matter in Light of Direct Detection Limits

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### Abstract

We examine the present and future XENON limits on the neutralino dark matter in split supersymmetry (split-SUSY). Through a scan over the parameter space under the current constraints from collider experiments and the WMAP measurement of the dark matter relic density, we find that in the allowed parameter space a large part has been excluded by the present XENON100 limits and a further largish part can be covered by the future exposure (6000 kg-day). In case of unobservation of dark matter with such an exposure in the future, the lightest neutralino will remain bino-like and its annihilation is mainly through exchanging the SM-like Higgs boson in order to get the required relic density.

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So far the only phenomenological crisis which requires new physics at the TeV scale seems to be the cosmic dark matter. Unlike neutrino oscillations which may indicate new physics at a very high inaccessible energy scale, the cosmic dark matter naturally points to a WIMP (weakly interacting massive particle) that should appear in some new physics around the TeV scale. A perfect candidate for such a WIMP is the lightest neutralino in low-energy supersymmetry (SUSY). As a specific low-energy SUSY model, split-SUSY is phenomenologically attractive because it just gives up the naturalness (fine-tuning) problem while maintaining the phenomenologically required dark matter and gauge coupling unification. This model also gets rid of the notorious supersymmetric flavor problem because of the assumed superheavy sfermions. In this framework, no scalar particles except the SM-like Higgs boson are accessible at foreseeable particle colliders. Consequently, the only way to explore this model is to study its gaugino/higgsino sector, for which dark matter detection experiments like XENON and CDMS can interplay with collider experiments to allow for a comprehensive test.

Recently, the CDMS and XENON collaborations reported their null search results which set rather stringent limits on the dark matter scattering cross section. The implications of these new limits for the neutralino dark matter in low-energy SUSY models have been discussed extensively. On the other hand, the CoGeNT and DAMA/LIBRA collaborations reported some excesses consistent with an explanation of light dark matter with a mass around 10 GeV (albeit not corroborated by CDMS or XENON results). The possible existence of such light dark matter has also stimulated theoretical studies in low-energy SUSY models.

In this note we discuss the implications of direct detection limits for the neutralino dark matter in split-SUSY. Since the most stringent limits come from the XENON100 results, we focus on the present and future (6000 kg-day) limits from XENON. We perform a scan over the parameter space under current constraints from collider experiments and the WMAP measurement of the dark matter relic density, and display the allowed parameter space in the plane of dark matter scattering rate versus dark matter mass. This allows us to see how much parameter space can be excluded by present and future XENON limits. Further, we show the implications of XENON limits on the properties of the neutralino dark matter and the lightest chargino.

We begin our analysis by writing the chargino mass matrix:

$$M_{\chi^\pm} = \begin{pmatrix} \sqrt{2} m_W \cos\beta & M_2 \\ \sqrt{2} m_W \sin\beta & \mu \end{pmatrix}$$

where the two-component spinors are defined as $\tilde{\psi}_+ = (i\tilde{\omega}^+, \tilde{h}_1^+)^T$ and $\tilde{\psi}_- = (i\tilde{\omega}^-, \tilde{h}_2^-)^T$. The neutralino mass matrix is given by:

$$M_{\chi^0} = \begin{pmatrix} M_1 & 0 & -m_Z \sin\theta_W \cos\beta & m_Z \sin\theta_W \sin\beta \\ 0 & M_2 & m_Z \cos\theta_W \cos\beta & -m_Z \cos\theta_W \sin\beta \\ -m_Z \sin\theta_W \cos\beta & m_Z \cos\theta_W \cos\beta & 0 & -\mu \\ m_Z \sin\theta_W \sin\beta & -m_Z \cos\theta_W \sin\beta & -\mu & 0 \end{pmatrix}$$

where the two-component spinors are defined as $\tilde{\psi}^0 = (i\tilde{\omega}^3, \tilde{h}_1, \tilde{h}_2)^T$. In these mass matrices, $M_1$ and $M_2$ are respectively the $U(1)$ and $SU(2)$ gaugino mass parameters, $\mu$ is the mass parameter in the mixing term $\mu\epsilon_{ij}H_u^i H_d^j$ in the superpotential, and $\tan\beta = v_2/v_1$ is the ratio of the vacuum expectation values of the two Higgs doublets.

The chargino mass matrix is diagonalized by $U^*$ and $V$ with the convention $M_{\chi^\pm} = U^* M_{\chi^\pm}^{\text{diag}} V^\dagger$ to give two chargino mass eigenstates $\chi_{1,2}^\pm$ with the convention $M_{\chi_1^\pm} < M_{\chi_2^\pm}$. The eigenstates may be wino-dominant or higgsino-dominant. Similarly, the neutralino mass matrix is diagonalized by $N^*$ with the convention $M_{\chi^0} = N^* M_{\chi^0}^{\text{diag}} N^\dagger$ to give four neutralino mass eigenstates $\chi_{1,2,3,4}^0$ with the convention $M_{\chi_1^0} < M_{\chi_2^0} < M_{\chi_3^0} < M_{\chi_4^0}$. The neutralinos may be bino-dominant, wino-dominant, or higgsino-dominant. Thus, the masses and mixings of charginos and neutralinos are determined by four parameters: $M_1$, $M_2$, $\mu$, and $\tan\beta$.

The spin-independent (SI) interaction between the lightest neutralino $\tilde{\chi}_1^0$ and the nucleon (denoted by $f_p$ for proton and $f_n$ for neutron) is induced by exchanging the SM-like Higgs boson or the squarks at tree level. In split-SUSY, the squark contribution is negligibly small, so $f_p$ is approximated by (similarly for $f_n$):

$$f_p \simeq \sum_{q=u,d,s} m_p f_{T_q}^{(p)} \frac{C_{h\tilde{\chi}\tilde{\chi}} C_{hqq}}{m_h^2} + \frac{2}{27} f_{TG}^{(p)} \sum_{q=c,b,t} m_p \frac{C_{h\tilde{\chi}\tilde{\chi}} C_{hqq}}{m_h^2}$$

where $f_{T_q}^{(p)}$ denotes the fraction of the proton mass from a light quark $q$ while $f_{TG}^{(p)} = 1 - \sum_{q=u,d,s} f_{T_q}^{(p)}$ is the heavy quark contribution through gluon exchange. $C_{h\tilde{\chi}\tilde{\chi}}$ is the coefficient of the effective scalar operator given by $C_{h\tilde{\chi}\tilde{\chi}} = C_{hqq}$ with $C$ standing for the corresponding Yukawa couplings. The $\tilde{\chi}^0$-nucleus scattering rate is then given by:

$$\sigma_{\text{SI}} = \frac{4}{\pi} \left( \frac{m_{\tilde{\chi}^0} m_T}{m_{\tilde{\chi}^0} + m_T} \right)^2 \left[ n_p f_p + n_n f_n \right]^2$$

where $m_T$ is the mass of the target nucleus and $n_p$ ($n_n$) is the number of protons (neutrons) in the target nucleus.

From these formulas we can infer when the scattering cross section is large. The expression indicates this occurs when $C_{h\tilde{\chi}\tilde{\chi}}$ and/or $C_{hqq}$ are enhanced. As the

Higgs boson is SM-like, $C_{hqq}$ has no $\tan\beta$ enhancement for down-type quarks. We therefore examine the behavior of $C_{h\tilde{\chi}\tilde{\chi}}$ with variation of the relevant SUSY parameters. For a bino-like $\tilde{\chi}_1^0$, this coupling is generated through bino-higgsino mixing, and thus a large $C_{h\tilde{\chi}\tilde{\chi}}$ requires large mixing, which means small $\mu$. To make this statement clearer, we consider the limit $M_1 \ll M_2, \mu$ (where $M_1$, $M_2$, and $\mu$ denote respectively the masses of bino, wino, and higgsino). After diagonalizing the neutralino mass matrix perturbatively, one obtains:

$$C_{h\tilde{\chi}\tilde{\chi}} \simeq \frac{m_Z \sin\theta_W \tan\theta_W}{\sqrt{2}} \frac{M_1 + \mu \sin 2\beta}{M_1 \mu}$$

Thus, the coupling $C_{h\tilde{\chi}\tilde{\chi}}$ becomes large when $\mu$ approaches $M_1$ from above.

In our numerical calculation for the dark matter-nucleon scattering rate, we considered all contributions (including QCD corrections) known to date. We take $f_{T_u}^{(p)} = 0.023$, $f_{T_d}^{(p)} = 0.032$, $f_{T_s}^{(p)} = 0.020$ and $f_{T_u}^{(n)} = 0.017$, $f_{T_d}^{(n)} = 0.041$, $f_{T_s}^{(n)} = 0.020$. Note that the value of $f_{T_s}$ is much smaller than that used in most previous studies. This small value comes from recent lattice simulations and can reduce the scattering rate significantly.

For the calculation of the SM-like Higgs boson mass, since in split-SUSY we have $\log(m_{\tilde{f}}^2/m_Z^2) \sim 1$ which spoils the convergence of the traditional loop expansion in evaluating SUSY effects on the Higgs boson self-energy, we use the effective potential method which involves the renormalization group evolution of SUSY effects from the squark scale to the electroweak scale. This computational method is employed in the package NMSSMTools. This package, which primarily serves as an important tool for studying the phenomenology of the Next-to-Minimal Supersymmetric Model, can also be applied to the MSSM case by setting $\lambda = \kappa \to 0$ (with this setting the singlet superfield decouples from the rest of the theory so that the MSSM phenomenology is recovered). Throughout our calculations we use this package.

As shown in previous work, the effects of sfermions and heavy Higgs bosons on electroweak theory begin to decouple when these particles are heavier than several TeV. Therefore, in our analysis we set $m_{\tilde{f}} = m_A = M_0 = 10$ TeV and the trilinear terms $A_t = A_b = 0$ TeV to simulate the split-SUSY scenario. We checked that the results with $M_0 = 100$ TeV are quite similar to those with $M_0 = 10$ TeV.

In the following we examine the properties of the parameter space surviving the present XENON experiment. As shown in the analytical expression, as $\mu$ approaches $M_1$, the coupling $C_{h\tilde{\chi}\tilde{\chi}}$ can be enhanced. This is reflected in the left panel of Fig. 2, where one can see that for most points excluded by the present XENON limits or covered by the future XENON exposure, they lie in the region where $M_1 \simeq \mu$ so that the $\tilde{\chi}_1^0$-nucleon scattering cross section is large. In contrast, the remaining inaccessible points fall into a region (denoted by green '×' symbols) where $\mu$ is much larger than $M_1$. This inaccessible region is shown again in the right panel of Fig. 2 [Figure 2: see original paper], which indicates that $m_{\chi_1^\pm}$ is about $2m_{\chi_1^0}$ in this region. The reason is that in this region the lightest neutralino is bino-like and the lightest chargino is wino-like. In the accessible region where $M_1 \simeq \mu$, the chargino has a large higgsino component (and a small wino component), while in the inaccessible region the chargino has a small higgsino component (and a large wino component).

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## References

[3] Z. Ahmed et al., Science 327, 1619 (2010).
[4] J. Hisano, K. Nakayama and M. Yamanaka, Phys. Lett. B 684, 246 (2010).
[5] M. Asano et al., arXiv:0912.5361; M. Holmes and B. D. Nelson, Phys. Rev. D 81, 055002 (2010).
[6] J. Cao et al., Phys. Rev. D 82, 051701 (2010); JHEP1007, 044 (2010).
[7] C. E. Aalseth et al., arXiv:1002.4703; arXiv:1106.0650; D. Hooper and C. Kelso, arXiv:1106.1066.
[8] R. Bernabei et al., Eur. Phys. J. C67, 39 (2010).
[9] D. A. Vasquez et al., Phys. Rev. D 82, 115027 (2010); P. Draper et al., Phys. Rev. Lett. 106, 121805 (2011); J. Cao et al., arXiv:1104.1754; D. Das and U. Ellwanger, JHEP 1009, 085 (2010); J. F. Gunion et al., arXiv:1009.2555; L. Calibbi, T. Ota, Y. Takanishi, arXiv:1104.1134; N. Fornengo et al., Phys. Rev. D 83, 015001 (2011); Z. Kang et al, JCAP 1101, 028 (2011).
[10] G. Jungman, M. Kamionkowski and K. Griest, Phys. Rept. 267, 195 (1996).
[11] M. Drees and M. Nojiri, Phys. Rev. D 48, 3483 (1993).
[12] A. Djouadi and M. Drees, Phys. Lett. B 484, 183 (2000); G. Belanger et al., Comput. Phys. Commun. 180, 747 (2009).
[13] M. S. Carena et al., Nucl. Phys. B 577, 88 (2000).
[14] J. Hisano, K. Ishiwata and N. Nagata, arXiv:1007.2601 [hep-ph].
[15] H. Ohki et al., Phys. Rev. D 78, 054502 (2008); D. Toussaint and W. Freeman, Phys. Rev. Lett. 103, 122002 (2009); J. Giedt, A. W. Thomas and R. D. Young, Phys. Rev. Lett. 103, 201802 (2009).
[16] M. Binger, Phys. Rev. D73, 095001 (2006). [hep-ph/0408240].
[17] U. Ellwanger et al., JHEP 0502, 066 (2005).
[18] U. Ellwanger, C. Hugonie, A. M. Teixeira, Phys. Rept. 496, 1-77 (2010). [arXiv:0910.1785 [hep-ph]].
[19] F. Wang, W. Wang and J. M. Yang, Eur. Phys. Jour. C 46, 521 (2006); Eur. Phys. Jour. C 51, 713 (2007).
[20] P. Gondolo et al., JCAP 07 (2004) 008. The code is available from http://www.physto.se/edsjo/darksusy
[21] J. Dunkley et al. [WMAP Collaboration], Astrophys. J. Suppl. 180, 306 (2009).
[22] J. Cao and J. M. Yang, JHEP 0812, 006 (2008).
[23] D. Feldman, K. Freese, P. Nath, B. D. Nelson and G. Peim, Phys. Rev. D 84, 015007 (2011).

[24] S. Akula, D. Feldman, P. Nath, G. Peim, arXiv:1107.3535.

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