

SLAC-PUB-11487
UCB-PTH-05/29
LBNL-58876

Dark Matter before the LHC in a Natural Supersymmetric Standard Model

Ryuichiro Kitano^a and Yasunori Nomura^{b,c}^a *Stanford Linear Accelerator Center, Stanford University, Stanford, CA 94309*^b *Department of Physics, University of California, Berkeley, CA 94720*^c *Theoretical Physics Group, Lawrence Berkeley National Laboratory, Berkeley, CA 94720*

Abstract

We show that the solid lower bound of about 10^{-44} cm² is obtained for the cross section between the supersymmetric dark matter and nucleon in a theory in which the supersymmetric fine-tuning problem is solved without extending the Higgs sector at the weak scale. This bound arises because of relatively small superparticle masses and a fortunate correlation that the two dominant diagrams for the dark matter detection always interfere constructively if the constraint from the $b \rightarrow s\gamma$ measurements is obeyed. It is, therefore, quite promising in the present scenario that the supersymmetric dark matter is discovered before the LHC, assuming that the dark matter is the lightest supersymmetric particle.

Weak scale supersymmetry provides an elegant framework to address the naturalness problem of the standard model as well as to provide the dark matter of the universe. On the other hand, there are already strong constraints on how supersymmetry can be realized at the weak scale. One of the severest constraints comes from the LEP II bound on the Higgs boson mass [1]. To evade this bound, the top squarks should generically be heavy, which in turn gives a large negative correction to the Higgs mass-squared parameter. This requires fine-tuning of order a few percent or worse to reproduce the correct scale for electroweak symmetry breaking in many supersymmetric models. Since our main motivation for weak scale supersymmetry comes from the concept of naturalness, we do not want fine-tuning in a theory with supersymmetry. This problem is called the supersymmetric fine-tuning problem. Addressing this issue may provide a key to find the correct theory at the weak scale, but it is generically difficult to solve the problem without extending the structure of the Higgs sector substantially.

Recently, we have shown that the supersymmetric fine-tuning problem can be solved within the framework of the minimal supersymmetric standard model (MSSM) [2]. Crucial ingredients for the solution are a large A term for the top squarks and a small B term for the Higgs doublets, which ensure successful electroweak symmetry breaking with relatively small top squark masses. A one-loop correction from the top-stop loop to the Higgs mass-squared parameter is made small by effectively lowering the messenger scale of supersymmetry breaking [3, 4], by invoking a mixture of moduli [5] and anomaly mediated [6] contributions as a source of supersymmetry breaking [7]. An interesting aspect of this solution is that all the superparticles are relatively light, with the masses determined only by a few free parameters. To evade fine-tuning worse than 20%, we find that the gauginos, whose masses are almost universal at the weak scale, must be lighter than about 900 GeV. The squarks and sleptons are even lighter, with the masses a factor of $\sqrt{2}$ smaller than that of the gauginos. Such light superparticles can have large impacts on the detections of these particles at colliders and in other experiments.

In this paper we show that the detection of supersymmetric dark matter is quite promising in the framework described above for solving the supersymmetric fine-tuning problem. In this framework, the lighter neutral Higgsino is a good candidate for the dark matter. First, it is lighter than about 200 GeV from the naturalness requirement, and thus is the lightest supersymmetric particle (LSP). Second, it can be generated nonthermally by the decay of the gravitino, which in turn is generated by the decay of the moduli field. For the range of the gravitino and moduli masses relevant to us, the right abundance of $\Omega_{\tilde{h}_0} \simeq 0.2$ can be naturally obtained [8]. A comparable amount of relic Higgsinos may also arise from the decay of Q-balls formed in the early stage of baryogenesis [9]. In any event, once we have the Higgsino as the dominant component of the dark matter, we can calculate its detection rates independently of the production mechanism.

We find that in the parameter region relevant to our solution to the fine-tuning problem,

the cross sections between the Higgsino and nuclei are dominated by the t -channel Higgs boson exchange diagrams. There are two contributions coming from the lighter and heavier Higgs boson exchanges, which are generically comparable in size. We find that the two contributions interfere constructively (destructively) if the sign of the μ parameter is positive (negative) in the standard convention. On the other hand, there is a strong constraint on the sign of μ from the measurements of the $b \rightarrow s\gamma$ process. We find that if the sign of μ is positive a large portion of the parameter space is consistent with the current measurements, while the case with negative μ is almost completely excluded. Because of this fortunate correlation, together with the fact that Higgsino dark matter generically has larger cross sections to nuclei than bino dark matter, the detection of supersymmetric dark matter in our framework is quite promising. We find that the parameter region relevant for the solution to the fine-tuning problem gives the dark matter-nucleon spin-independent cross section ranging from about 10^{-44} cm² to several times 10^{-42} cm². Some of the parameter region, therefore, is even already excluded by the latest data from the CDMS II experiment [10]. Since we obtain the lower bound on the cross section of about 10^{-44} cm², we expect that most of the relevant parameter space will be covered by further running of the CDMS II in the next two years.

Let us now start calculating the cross section between the Higgsino dark matter and nuclei in the present framework. For this purpose, we need the masses and interactions of the superparticles and the Higgs bosons. The soft supersymmetry breaking parameters for the gauginos and sfermions are given by

$$M_1 = M_2 = M_3 = M_0, \quad (1)$$

$$m_{\tilde{q}}^2 = m_{\tilde{u}}^2 = m_{\tilde{d}}^2 = m_{\tilde{l}}^2 = m_{\tilde{e}}^2 = \frac{M_0^2}{2}, \quad (2)$$

$$A_u = A_d = A_e = M_0, \quad (3)$$

at the effective messenger scale $M_{\text{mess}} \sim \text{TeV}$, where the parameter M_0 takes a value in the range

$$450 \text{ GeV} \lesssim M_0 \lesssim 900 \text{ GeV}, \quad (4)$$

for $\tan\beta \equiv \langle H_u \rangle / \langle H_d \rangle \gtrsim 20$. For smaller $\tan\beta$, the lower bound in Eq. (4) increases; for $\tan\beta = 10$ (5) the lower bound becomes ≈ 550 GeV (900 GeV). The range of $\tan\beta$ we consider is thus

$$\tan\beta \gtrsim 5. \quad (5)$$

The upper bound in Eq. (4) comes from requiring the absence of fine-tuning worse than 20%: $\Delta^{-1} \geq 20\%$. The Higgs sector of the theory contains four free parameters, $m_{H_u}^2$, $m_{H_d}^2$, μ and B , but one combination is fixed by the vacuum expectation value $v \equiv (\langle H_u \rangle^2 + \langle H_d \rangle^2)^{1/2} \simeq 174$ GeV, leaving only three independent parameters. We take these parameters to be μ , m_A

and $\tan \beta$, where m_A is the mass of the pseudo-scalar Higgs boson. The naturalness requirement, $\Delta^{-1} \geq 20\%$, then leads to the bounds

$$|\mu| \lesssim 190 \text{ GeV}, \quad (6)$$

and

$$m_A \lesssim 300 \text{ GeV}. \quad (7)$$

Overall, the number of free parameters in our analysis is five: M_0 , M_{mess} , $\tan \beta$, μ and m_A . The values of M_0 , $\tan \beta$, μ and m_A are bounded by Eqs. (4), (5), (6) and (7), respectively, and M_{mess} has to be close to TeV: $50 \text{ GeV} \lesssim M_{\text{mess}} \lesssim 3 \text{ TeV}$.

From Eqs. (1 – 4) and (6), we find that the Higgsinos, which have masses of about $|\mu|$, are significantly lighter than the other superparticles, so that the LSP comes from one of the Higgsinos. After electroweak symmetry breaking, there are three Higgsino states: two neutral states \tilde{h}^0 and \tilde{h}'^0 , defined here by $m_{\tilde{h}'^0} > m_{\tilde{h}^0}$, and one charged state \tilde{h}^\pm . The mass splittings between these states are approximately given by

$$m_{\tilde{h}'^0} - m_{\tilde{h}^\pm} \simeq m_{\tilde{h}^\pm} - m_{\tilde{h}^0} \simeq \frac{m_Z^2}{2M_0}, \quad (8)$$

where the corrections are only of order 20% in the relevant parameter region. Therefore, we find that the lighter neutral Higgsino is always the LSP, and the mass splittings between the different Higgsino states are of order 10 GeV.

Suppose now that the neutral Higgsino \tilde{h}^0 , produced nonthermally in the early universe, is the dominant component of the dark matter. Since the mass splittings in Eq. (8) are much larger than the kinetic energy of the dark matter, we only have to consider elastic scattering between \tilde{h}^0 and nuclei to study the detection of the Higgsino dark matter. The scattering occurs through exchanges of the squarks and Higgs bosons, but in the parameter region relevant to us the Higgs boson exchange diagrams dominate. We thus focus on the Higgs boson exchange contributions below. The squark exchange contributions, however, are also included in our numerical analysis given later.

The exchange of the Higgs bosons induces the following effective interactions between the dark matter χ and the quarks $q = u, d, s, c, b, t$ [11]:

$$\mathcal{L}_{\text{eff}} = \sum_q f_q m_q \bar{\chi} \chi \bar{q} q, \quad (9)$$

where m_q are the quark masses and f_q are given by

$$f_q \simeq -\frac{g^2 + g'^2}{8(M_0 - |\mu|)} \left(\frac{c_{h\chi\chi} c_{hqq}}{m_h^2} + \frac{c_{H\chi\chi} c_{Hqq}}{m_H^2} \right). \quad (10)$$

Here, m_h and m_H are the masses of the lighter and heavier neutral Higgs bosons, h and H , and the coefficients c 's are given by

$$c_{h\chi\chi} \simeq 1 + \text{sgn}(\mu) \sin 2\beta, \quad c_{H\chi\chi} \simeq -\text{sgn}(\mu) \cos 2\beta, \quad (11)$$

$$c_{hq\bar{q}} = \begin{cases} \frac{\cos \alpha}{\sin \beta} & (\text{for } q = u, c, t) \\ -\frac{\sin \alpha}{\cos \beta} & (\text{for } q = d, s, b) \end{cases}, \quad c_{Hq\bar{q}} = \begin{cases} \frac{\sin \alpha}{\sin \beta} & (\text{for } q = u, c, t) \\ \frac{\cos \alpha}{\cos \beta} & (\text{for } q = d, s, b) \end{cases}, \quad (12)$$

where α is the mixing angle for the neutral Higgs bosons ($-\pi/2 \leq \alpha \leq 0$). The spin-independent cross section of the dark matter with a target nucleus is then given by

$$\sigma_T = \frac{4}{\pi} \left(\frac{m_\chi m_T}{m_\chi + m_T} \right)^2 (n_p f_p + n_n f_n)^2, \quad (13)$$

where m_T is the mass of the target nucleus, and n_p and n_n are the proton and neutron numbers in the nucleus, respectively. The quantities f_p and f_n are given by

$$f_N = m_N \sum_q f_q \mathcal{M}_q^{(N)}, \quad (14)$$

where $N = p, n$. Here, $\mathcal{M}_q^{(N)}$ are the nucleon matrix elements defined by $\langle N | m_q \bar{q}q | N \rangle = m_N \mathcal{M}_q^{(N)}$, whose approximate values are given by

$$\mathcal{M}_u^{(p)} \simeq 0.023, \quad \mathcal{M}_d^{(p)} \simeq 0.034, \quad \mathcal{M}_s^{(p)} \simeq 0.14, \quad \mathcal{M}_c^{(p)} \simeq \mathcal{M}_b^{(p)} \simeq \mathcal{M}_t^{(p)} \simeq 0.059, \quad (15)$$

$$\mathcal{M}_u^{(n)} \simeq 0.019, \quad \mathcal{M}_d^{(n)} \simeq 0.041, \quad \mathcal{M}_s^{(n)} \simeq 0.14, \quad \mathcal{M}_c^{(n)} \simeq \mathcal{M}_b^{(n)} \simeq \mathcal{M}_t^{(n)} \simeq 0.059. \quad (16)$$

From these expressions, we find the following. First, Eq. (11) tells us that the signs of $c_{h\chi\chi}$ and $c_{H\chi\chi}$ are given by $\text{sgn}(c_{h\chi\chi}) = +1$ and $\text{sgn}(c_{H\chi\chi}) = \text{sgn}(\mu)$, respectively. On the other hand, from Eqs. (14 – 16), we find that the total contributions to the cross section from operators with down-type quarks, $q = d, s, b$, are larger than those from operators with up-type quarks, $q = u, c, t$. Combining these with the fact that $\text{sgn}(c_{hq\bar{q}}) = \text{sgn}(c_{Hq\bar{q}}) = +1$ for $q = d, s, b$ (see Eq. (12)), we find from Eq. (10) that the contributions from the light and heavy Higgs bosons interfere constructively (destructively) if the sign of μ is positive (negative). This implies that for $\mu > 0$ we have a solid lower bound on the cross section because our mass parameters M_0 , μ , m_h , and m_H are bounded by Eq. (4), Eq. (6), $m_h \lesssim 120$ GeV, and $m_H \simeq m_A \lesssim 300$ GeV (see Eq. (7)), respectively. Here, the bound on m_h follows from that on M_0 . The case with $\mu < 0$, on the other hand, admits a possibility of cancellation between the two contributions. Since $\alpha \sim \beta - \pi/2$ in our parameter region, an excessive cancellation occurs if $m_H^2/m_h^2 \sim \tan \beta$ (see Eqs. (10, 11, 12)).

In Fig. 1, we have plotted the cross section of the dark matter χ with a target nucleus X in the case of $X = \text{Ge}$, for $\mu > 0$ (left) and $\mu < 0$ (right). Here, we have plotted the cross section

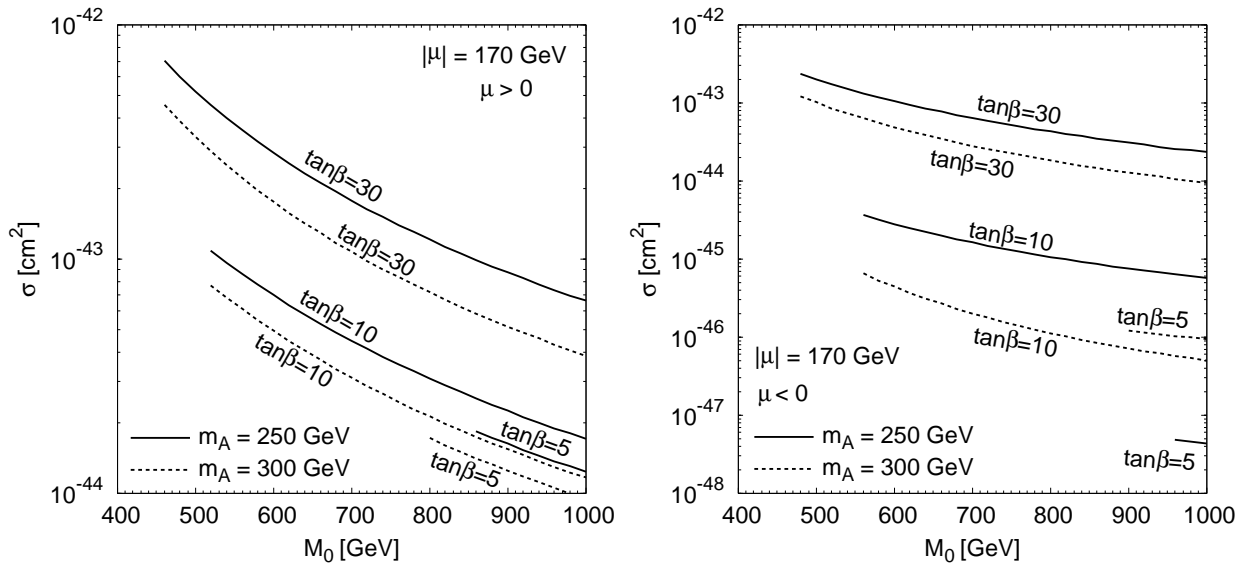


Figure 1: Dark matter-nucleon cross section, σ , as a function of M_0 for the case of $X = \text{Ge}$. The left and right plots correspond to the cases with $\mu > 0$ and $\mu < 0$, respectively. The parameters $|\mu|$ and M_{mess} are fixed as $|\mu| = 170 \text{ GeV}$ and $M_{\text{mess}} = 1 \text{ TeV}$, and $\tan \beta$ and m_A are varied as $\tan \beta = 5, 10, 30$ and $m_A = \{250, 300\} \text{ GeV}$. The curves are drawn only for the values of M_0 for which the Higgs boson mass bound of $m_h \gtrsim 114.4 \text{ GeV}$ is satisfied.

normalized to the nucleon [12]

$$\sigma = \frac{\sigma_T}{(n_p + n_n)^2} \left(\frac{1 \text{ GeV}}{m_\chi m_T / (m_\chi + m_T)} \right)^2, \quad (17)$$

as a function of M_0 for several values of $\tan \beta$ and m_A : $\tan \beta = 5, 10, 30$ and $m_A = \{250, 300\} \text{ GeV}$. In the figure, all the contributions discussed in [11] including the squark exchange contributions, as well as the effects of running from M_{mess} to the superparticle mass scale, are included. The curves are drawn only for the values of M_0 for which the LEP II bound on the Higgs boson mass, $m_h \gtrsim 114.4 \text{ GeV}$ [1], is satisfied. The parameters $|\mu|$ and M_{mess} are fixed as $|\mu| = 170 \text{ GeV}$ and $M_{\text{mess}} = 1 \text{ TeV}$, but the sensitivities to these parameters are not significant.

In the figure, we can clearly see the general features discussed before. First, the cross section is larger for $\mu > 0$ than $\mu < 0$. Second, we obtain the lower bound on σ of about 10^{-44} cm^2 for $\mu > 0$, since M_0 is bounded as $M_0 \lesssim 900 \text{ GeV}$ from the naturalness requirement. This is quite interesting because this region will be probed in the further running of the CDMS II experiment in the next two years. On the other hand, we find that the cross section can be as small as $(10^{-47} \sim 10^{-46}) \text{ cm}^2$ for $\mu < 0$. Therefore, fully exploring this case would require next generation experiments.

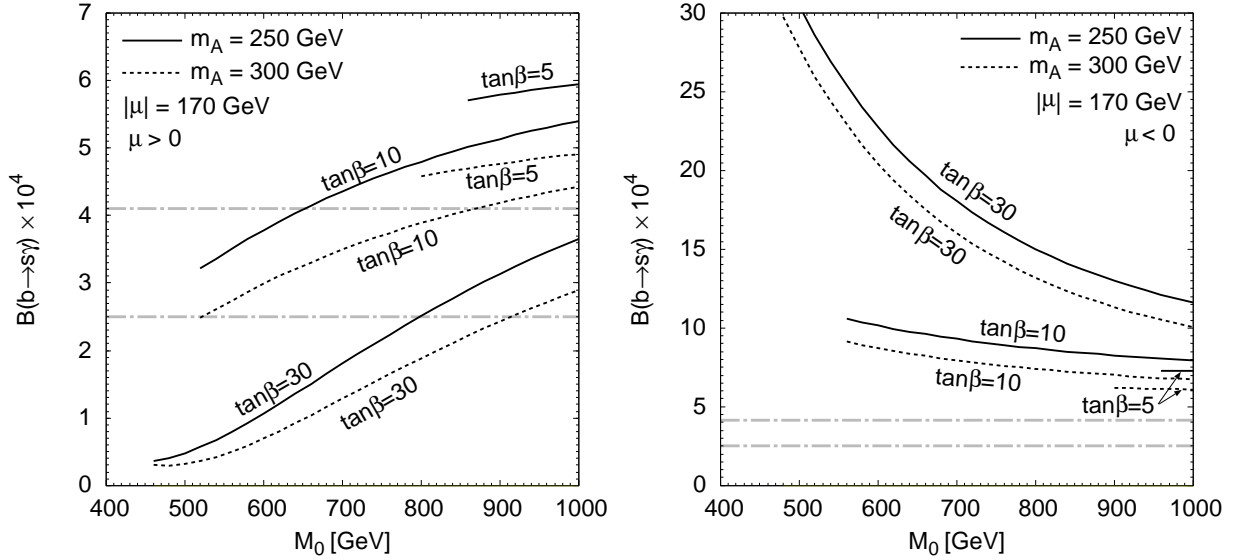


Figure 2: Branching ratio for the $b \rightarrow s\gamma$ process, $B(b \rightarrow s\gamma)$, as a function of M_0 . The left and right plots correspond to the cases with $\mu > 0$ and $\mu < 0$, respectively. The parameters $|\mu|$ and M_{mess} are fixed as $|\mu| = 170$ GeV and $M_{\text{mess}} = 1$ TeV, and $\tan\beta$ and m_A are varied as $\tan\beta = 5, 10, 30$ and $m_A = \{250, 300\}$ GeV. The curves are drawn only for the values of M_0 for which the Higgs boson mass bound of $m_h \gtrsim 114.4$ GeV is satisfied.

Can we say something about the sign of μ from other experiments? It is known that the rate for the $b \rightarrow s\gamma$ process can depend highly on the signs of μ and A_t . In our parameter region, contributions from chargino and charged Higgs boson loops interfere destructively (constructively) for $\mu > 0$ (< 0) in the $b \rightarrow s\gamma$ amplitude, implying that the positive sign of μ is preferred over the other one. This situation is depicted in Fig. 2, where we have plotted the branching ratio $B(b \rightarrow s\gamma)$ for $\mu > 0$ (left) and $\mu < 0$ (right), calculated using the DarkSUSY package [13]. We have again drawn the curves as a function of M_0 for $\tan\beta = 5, 10, 30$, $m_A = \{250, 300\}$ GeV, $|\mu| = 170$ GeV and $M_{\text{mess}} = 1$ TeV only for the regions where the bound on the Higgs boson mass is satisfied. The sensitivities to the fixed parameters, μ and M_{mess} , are weak in our parameter region. We have also indicated the 2σ allowed region from the current measurements, $B(b \rightarrow s\gamma) = (3.3 \pm 0.4) \times 10^{-4}$ [14], by two horizontal dot-dashed lines.

While the predictions for the rate of the $b \rightarrow s\gamma$ process are still subject to theoretical uncertainties of order (20~30)% [15], we can still conclude from Fig. 2 that the case with $\mu < 0$ is almost entirely excluded. On the other hand, for $\mu > 0$ a large portion of the parameter space is consistent with the data within theoretical and experimental uncertainties. The positive sign of μ may also be preferred from the measurement of the muon anomalous magnetic moment [16],

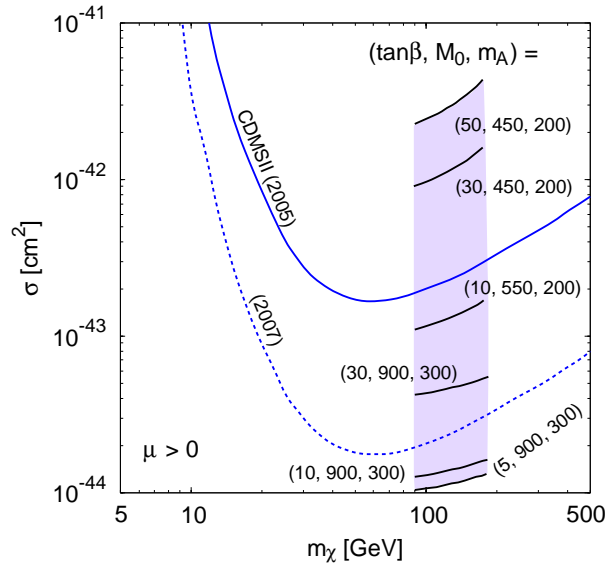


Figure 3: Dark matter-nucleon cross section, σ , as a function of the dark matter mass m_χ for various values of $(\tan\beta, M_0, m_A)$. The shaded area corresponds to our parameter region in which the supersymmetric fine-tuning problem is solved. The exclusion curve from the latest CDMS II data is also shown (a solid line), as well as an expected future sensitivity (a dashed line).

because supersymmetric corrections decrease (increase) the possible discrepancy between the standard model prediction and the data for $\mu > 0$ (< 0).

As we have seen before, the detection of our Higgsino dark matter is quite promising for $\mu > 0$. In Fig. 3, we have plotted the spin-independent cross section between the dark matter and nucleon as a function of the dark matter mass m_χ for various values of $\tan\beta$, M_0 and m_A . For $\tan\beta = 10$ (30), the values of M_0 and m_A are chosen such that they maximize, $(M_0, m_A) = (550, 200)$ ($(M_0, m_A) = (450, 200)$), or minimize, $(M_0, m_A) = (900, 300)$ ($(M_0, m_A) = (900, 300)$), the cross section, so that we can read off the allowed range of the cross section for a given value of $\tan\beta$ and m_χ . (The lower bound on m_A has been set to 200 GeV somewhat arbitrarily, motivated by the $b \rightarrow s\gamma$ data, but lowering this only extends the range of σ to the direction of larger σ .) The lines of $(\tan\beta, M_0, m_A) = (50, 450, 200)$ and $(5, 900, 300)$ have been drawn to set the upper and lower boundaries for the range of σ . Together with the experimental lower bound on m_χ and the naturalness bound on μ in Eq. (6), we obtain our predictions for σ and m_χ given by the shaded area in the figure. We have also drawn the exclusion curve from the latest CDMS II data [10] by a solid line, and an estimate for the expected future sensitivity (an order of magnitude improvement of the current bound by the end of 2007 [17]) by a dashed line.

It is interesting that the latest data from the CDMS II already exclude a portion of parameter

space relevant for our solution to the supersymmetric fine-tuning problem. For $\tan\beta = 30$, for example, the region $M_0 \lesssim 750$ GeV (500 GeV) is already excluded for $m_A = 200$ GeV (300 GeV). The current data are not yet very sensitive to smaller values for $\tan\beta$. A large portion of the relevant parameter space, however, will be covered by the end of 2007 by a planned order of magnitude improvement of the current lower bound on σ . For $\tan\beta = 30$, for example, this covers the entire parameter region relevant for the solution to the supersymmetric fine-tuning problem. Even for $\tan\beta = 10$, the entire region of M_0 is covered for $m_A = 200$ GeV, and $M_0 \lesssim 700$ GeV is covered for $m_A = 300$ GeV. It is, therefore, quite promising in the present framework that the supersymmetric dark matter is discovered before the LHC, assuming that the dominant component of the dark matter is in fact the lightest supersymmetric particle.

In summary, we have shown that the detection of supersymmetric dark matter is quite promising in the theory in which the supersymmetric fine-tuning problem is solved. In fact, this conclusion is somewhat more generic in the sense that in a theory where the supersymmetric fine-tuning is reduced, the Higgsino is generically light and is naturally a candidate for the LSP. Then, if the gaugino masses are not very large, as suggested indirectly by the fine-tuning argument, the detection rate is generically not very small. To obtain the solid lower bound on the detection rate, however, we need more information. We have shown that the $b \rightarrow s\gamma$ process can provide such information in our particular context, and that we can obtain the lower bound on the dark matter-nucleon cross section, $\sigma \gtrsim 10^{-44}$ cm², which is close to the expected sensitivity of the CDMS II experiment within the next two years. We find it quite interesting that the supersymmetric dark matter may in fact be discovered before the LHC turns on.

Acknowledgments

Y.N. thanks Jeffrey Filippini for conversation. The work of R.K. was supported by the U.S. Department of Energy under contract number DE-AC02-76SF00515. The work of Y.N. was supported in part by the Director, Office of Science, Office of High Energy and Nuclear Physics, of the US Department of Energy under Contract DE-AC02-05CH11231, by the National Science Foundation under grant PHY-0403380, by a DOE Outstanding Junior Investigator award, and by an Alfred P. Sloan Research Fellowship.

References

- [1] R. Barate *et al.* [ALEPH Collaboration], Phys. Lett. B **565**, 61 (2003) [arXiv:hep-ex/0306033]; LEP Higgs Working Group Collaboration, arXiv:hep-ex/0107030.
- [2] R. Kitano and Y. Nomura, arXiv:hep-ph/0509039.
- [3] K. Choi, K. S. Jeong and K. i. Okumura, arXiv:hep-ph/0504037.
- [4] K. Choi, K. S. Jeong, T. Kobayashi and K. i. Okumura, arXiv:hep-ph/0508029.
- [5] V. S. Kaplunovsky and J. Louis, Phys. Lett. B **306**, 269 (1993) [arXiv:hep-th/9303040]; A. Brignole, L. E. Ibanez and C. Munoz, Nucl. Phys. B **422**, 125 (1994) [Erratum-ibid. B **436**, 747 (1995)] [arXiv:hep-ph/9308271].
- [6] L. Randall and R. Sundrum, Nucl. Phys. B **557**, 79 (1999) [arXiv:hep-th/9810155]; G. F. Giudice, M. A. Luty, H. Murayama and R. Rattazzi, JHEP **9812**, 027 (1998) [arXiv:hep-ph/9810442].
- [7] For recent study of this scenario, see, e.g., K. Choi, A. Falkowski, H. P. Nilles, M. Olechowski and S. Pokorski, JHEP **0411**, 076 (2004) [arXiv:hep-th/0411066]; K. Choi, A. Falkowski, H. P. Nilles and M. Olechowski, Nucl. Phys. B **718**, 113 (2005) [arXiv:hep-th/0503216]; M. Endo, M. Yamaguchi and K. Yoshioka, Phys. Rev. D **72**, 015004 (2005) [arXiv:hep-ph/0504036]; A. Falkowski, O. Lebedev and Y. Mambrini, arXiv:hep-ph/0507110.
- [8] K. Kohri, M. Yamaguchi and J. Yokoyama, arXiv:hep-ph/0502211.
- [9] M. Fujii and K. Hamaguchi, Phys. Rev. D **66**, 083501 (2002) [arXiv:hep-ph/0205044].
- [10] D. S. Akerib *et al.* [CDMS Collaboration], arXiv:astro-ph/0509259.
- [11] M. Drees and M. Nojiri, Phys. Rev. D **48**, 3483 (1993) [arXiv:hep-ph/9307208], and references therein.
- [12] J. D. Lewin and P. F. Smith, Astropart. Phys. **6**, 87 (1996).
- [13] P. Gondolo, J. Edsjo, P. Ullio, L. Bergstrom, M. Schelke and E. A. Baltz, JCAP **0407**, 008 (2004) [arXiv:astro-ph/0406204].
- [14] S. Eidelman *et al.* [Particle Data Group Collaboration], Phys. Lett. B **592**, 1 (2004).
- [15] See, e.g., P. Gambino and M. Misiak, Nucl. Phys. B **611**, 338 (2001) [arXiv:hep-ph/0104034]; M. Neubert, Eur. Phys. J. C **40**, 165 (2005) [arXiv:hep-ph/0408179].
- [16] G. W. Bennett *et al.* [Muon g-2 Collaboration], Phys. Rev. Lett. **92**, 161802 (2004) [arXiv:hep-ex/0401008].
- [17] R. Schnee, talk at *INT Underground Science Workshop*, University of Washington, Seattle, June–July 2005; see also R. Gaitskell and V. Mandic, <http://dmtools.brown.edu>