Constraints of dark matter direct detection experiments on the MSSM and implications for LHC Higgs searches

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Abstract

Assuming the lightest neutralino solely composes the cosmic dark matter, we examine the constraints of the CDMS-II and XENON100 dark matter direct searches on the parameter space of the MSSM Higgs sector. We find that the current CDMS-II/XENON100 limits can exclude some of the parameter space which survive the constraints from the dark matter relic density and various collider experiments. We also find that in the currently allowed parameter space, the charged Higgs boson is hardly accessible at the LHC for an integrated luminosity of 30 fb\(^{-1}\), while the neutral non-SM Higgs bosons (\(H, A\)) may be accessible in some allowed region characterized by a large \(\mu\). The future XENON100 (6000 kg-days exposure) will significantly tighten the parameter space in case of nonobservation of dark matter, further shrinking the likelihood of discovering the non-SM Higgs bosons at the LHC.

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Introduction: The existence of non-baryonic cold dark matter (DM) has been established by cosmological observations [1]. Weakly interacting massive particles (WIMPs) are the natural candidates of DM, among which the lightest neutralino $\chi^0$ in the minimal supersymmetric standard model (MSSM) has been most extensively studied [2].

The most convincing detection for the neutralino DM is the underground direct detection experiments like CDMS and XENON, which search for neutralino-nucleon ($\chi N$) scattering in a low-background circumstance [3, 4]. Recently, both CDMS-II and XENON100 reported their search results [3, 4], which immediately stimulated some theoretical works [5, 6]. On the theoretical side, great efforts have been paid to improve the accuracy of the prediction for the $\chi N$ scattering rate [7, 8]. For example, it has long been known that the hadronic uncertainty, especially the strange quark content in a nucleon, can affect the rate by almost one order of magnitude, and is therefore impacting significantly the interpretation of the experimental searches for DM [9]. This problem was recently better solved by lattice simulation and it is found that the strange quark content is much smaller than previously thought, which leads to a significant suppression on the uncertainty [10].

In light of the above experimental and theoretical progress in DM study, we in this work re-investigate the $\chi N$ scattering and use the recent CDMS-II/XENON limits to constrain the MSSM parameter space. Note that unlike most recent studies which try to explain the two possible DM events reported by CDMS-II, we use the CDMS-II 90% upper limit on the spin-independent (SI) $\chi N$ scattering cross section. In addition to the direct detection limits from CDMS-II and XENON100, we also consider the constraints from the DM relic density and various collider experiments. We will first perform a scan over the MSSM parameter space by considering these constraints. Then we investigate the $\chi N$ scattering in the surviving parameter space to demonstrate the further constraints of CDMS-II/XENON on it. Given the extreme importance of Higgs search at the LHC and the strong correlation between the $\chi N$ scattering and the Higgs sector, our study will be focused on the MSSM Higgs sector.

$\chi N$ scattering in the MSSM: For the sensitivity in current DM direct detection experiments, it is sufficient to consider only the SI interactions between $\chi^0$ and nucleon (denoted by $f_p$ for proton and $f_n$ for neutron [11]) in calculating the scattering rate. In the MSSM, these interactions are induced by exchanging squarks or neutral Higgs bosons at tree level [7, 11]. For moderately light Higgs bosons, the latter contribution is dominant and $f_p$ is...
approximated by [11] (similarly for $f_n$)

$$f_p \simeq \sum_{q=u,d,s} \frac{f^H_q}{m_q} m_p f_T^{(p)} + \frac{2}{27} f_T \sum_{q=c,b,t} \frac{f^H_q}{m_q} m_p,$$

(1)

where $f_T^{(p)}$ denotes the fraction of $m_p$ (proton mass) from the light quark $q$ while $f_T = 1 - \sum_{u,d,s} f_T^{(p)}$ is the heavy quark contribution through gluon exchange. $f^H_q$ is the coefficient of the effective scalar operator given by [11]

$$f^H_q = m_q g^2 \left( \frac{C_{HXX} C_{Hqq}}{m_h^2} + \frac{C_{HXX} C_{Hqq}}{m_H^2} \right),$$

(2)

with $C$ standing for the corresponding Yukawa couplings. The $\chi^0_1$–nucleus scattering rate is then given by [11]

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

(3)

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

From the above formulas we can infer in which situation the scattering cross section is large. Eq.(2) indicates that this occurs only when $C_{SXX}$ and/or $C_{Sqq}$ ($S$ stands for a Higgs boson) get enhanced. Since the potential enhancement of $C_{Hdd}$ by $\tan \beta$ is well known, we here only analysis the behavior of $C_{SXX}$ with the variation of SUSY parameters. For a bino-like $\chi^0_1$ encountered in this work, this coupling is generated through the bino-higgsino mixing [2], so a large $C_{SXX}$ needs a large mixing, which means a small $\mu$. To make this statement clearer, one may consider the limit of $M_1 \ll M_2, \mu$ with $M_1, M_2$ and $\mu$ denoting respectively the masses of bino, wino and higgsino. After diagonalizing the neutralino mass matrix perturbatively, one can get [5]

$$C_{bXX} \simeq \frac{m_Z \sin \theta_W \tan \theta_W}{M_1^2 - \mu^2} \left[ M_1 + \mu \sin 2\beta \right],$$

$$C_{HXX} \simeq -\frac{m_Z \sin \theta_W \tan \theta_W}{M_1^2 - \mu^2} \mu \cos 2\beta.$$

(4)

So both couplings become large when $\mu$ approaches downward to $M_1$.

In our numerical calculations for the scattering rate, we considered all the contributions known so far, including the QCD correction, SUSY-QCD correction [8] as well as the contribution from high dimensional operators [11]. Note that the SUSY-QCD corrections are not negligible because they may sizably reduce the scattering rate by suppressing $C_{Sqq}$ [8, 12].
In our calculations we take $f_{T_u}^{(p)} = 0.023$, $f_{T_u}^{(n)} = 0.034$, $f_{T_u}^{(n)} = 0.019$, $f_{T_u}^{(n)} = 0.041$ and $f_{T_u}^{(p)} = f_{T_d}^{(n)} = 0.020$. Note the value of $f_{T_u}$ we choose is much smaller than that taken in most previous studies. This small value comes from the recent lattice simulation [10], and it can reduce the scattering rate significantly.

**Numerical results:** We make some assumptions to reduce the number of free parameters before our scan. First, we note that the first two generation squarks may be heavier than about 400 GeV from the Tevatron experiments [13], and thus their effects on the scattering should be unimportant in the presence of light Higgs bosons. So, for the first two generation squarks we fix the soft masses and the trilinear parameters to be 1 TeV. We checked our conclusion are not affected by such specific choice. Second, since the third generation squarks affect the Higgs sector significantly, we let free all the relevant soft parameters. But to simplify our analysis, we assume $m_{D_3} = m_{U_3}$ and $A_b = A_t$, which is well motivated by the mSUGRA model with large $\tan \beta$. Third, we note that, although the slepton masses do not directly affect the $\chi N$ scattering rate, they can affect the allowed range of $\tan \beta$ via the muon $g - 2$. In order to avoid a tight constraint on $\tan \beta$, we assume a universal soft parameter $m_{\tilde{\ell}}$ and vary it in our scan. Finally, we use the grand unification relation $3M_1/5\alpha_1 = M_2/\alpha_2 = M_3/\alpha_3$ for the gaugino masses.

With the above assumptions, the free parameters remained are scanned in the ranges: $1 \leq \tan \beta \leq 80$, $80 \text{ GeV} \leq m_A \leq 300 \text{ GeV}$, $30 \text{ GeV} \leq M_1 \leq 500 \text{ GeV}$, $100 \text{ GeV} \leq \mu, m_{\tilde{\ell}}, m_{\tilde{\chi}_3}, m_{\tilde{U}_3} \leq 1 \text{ TeV}$ and $-3 \text{ TeV} < A_t \leq 3 \text{ TeV}$. In our scan, we consider the following constraints like done in [14]: (1) Direct bounds on sparticle and Higgs masses from LEP and Tevatron experiments. (2) LEP II search for Higgs bosons, which includes various channels of Higgs boson productions. (3) LEP I and LEP II constraints on the productions of neutralinos and charginos. (4) Constraints ($2\sigma$) from precision electroweak observables plus $R_b$ [16], and also the constraints from B-physics observables such as $B \to X_s\gamma$, $B_s \to \mu^+\mu^-$, $B^+ \to \tau^+\nu$, and the mass difference $\Delta M_d$ and $\Delta M_s$. (5) The muon $g - 2$ constraint [17] (we require the MSSM contribution to explain the deviation at $2\sigma$ level). (6) We require $\chi^0_1$ to account for the WMAP measured dark matter relic density at $2\sigma$ level [1]. The samples surviving the above constraints will be input for the calculation of the $\chi N$ scattering rate. Note that most of the above constraints have been encoded in NMSSMTools [15]. We extended the code to the MSSM case, especially we wrote the code for the $\chi N$ scattering rate to improve the scan efficiency.
FIG. 1: The scatter plots for the spin-independent elastic cross section of $\chi N$ scattering under the constraints of dark matter relic density (2$\sigma$) and various collider experiments. The ‘+’ points (red) are excluded by CDMS II and XENON100 (90% C.L.) limits, the ‘×’ (blue) would be further excluded by XENON100 (6000 kg-days) in case of nonobservation, and the ‘◦’ (green) are beyond the XENON100 (6000 kg days) sensitivity.

To show the sensitivity of the $\chi N$ scattering rate to the value of $f_{Ts}$, we plot the surviving samples on the plane of the scattering rate versus the neutralino mass with the new lattice value $f_{Ts} = 0.02$ and with the old value $f_{Ts} = 0.38$ (corresponding to $\Sigma_{\pi N} \simeq 64$ MeV in [9]). One can see that the new lattice value of $f_{Ts}$ gives a much lower scattering rate. In our following results we fixed $f_{Ts} = 0.02$.

Our scan samples are $2 \times 10^{11}$ random points over the parameter space, and after considering the constraints, about $6 \times 10^{7}$ samples can survive. We find that for some survived samples the $\chi N$ scattering rate can be as large as $10^{-42}$ cm$^2$, which is far above the current CDMS II/XENON limits [3]. Requiring the scattering rate not exceed the current CDMS II/XENON bounds, we find that about 33% (61%) of the survived samples are ruled out for $m_A \leq 300$ GeV (200 GeV). Further, if the future XENON100 with 6000 kg-days exposure [18] gives null dark matter results, then 90.5% (99.5%) of the survived samples will be ruled out for $m_A \leq 300$ GeV (200 GeV). So the dark matter direct detection experiments are highly complementary to collider experiments in testing the MSSM.

In Fig. 2 the surviving samples are projected for $\tan \beta$ and $\mu$ versus the charge Higgs boson mass $m_{H^+}$. We see that the regions excluded by the CDMS-II/XENON limits are characterized by large $\tan \beta$ and small $\mu$. The future XENON100 (6000 kg-days) can further
shrink the currently allowed regions, and in particular set a bound $m_{H^+} \gtrsim 165$ GeV. Around this lower bound, the value of $\mu$ is quite large ($\simeq 1$ TeV) and so the MSSM has a fine-tuning due to the relation $m_{H^+}^2 > m_A^2$ and $m_A^2 = m_{h_u}^2 + m_{h_d}^2 + 2\mu^2$ [2].

In our following discussions, we only focus on the samples that satisfy the current CDMS-II/XENON limits. In Fig. 3 we project the surviving samples in the planes of $\tan \beta - \mu$ and $\tan \beta_{\text{eff}} - m_{H^+}$. Here $\tan \beta_{\text{eff}} \equiv \tan \beta / (1 + \Delta_b)$ with $\Delta_b$ denoting the SUSY radiative corrections to bottom quark mass[12]. As expected, large $\tan \beta$ must be accompanied by large $\mu$ to suppress the the scattering rate, and this tendency becomes more apparent for the samples further satisfying the future XENON100 limit. We note that in this case, i.e. large $\tan \beta$ along with large $\mu$, $\Delta_b$ should be large[12] so that $\tan \beta_{\text{eff}}$ is significantly smaller than $\tan \beta$. This speculation is verified by Fig. 3 and also by our results for $\Delta_b$, which show $\Delta_b$ larger than 30% for $\tan \beta \geq 40$.

Implication for LHC Higgs searches: Above results showed that the CDMS-II/XENON limits have set upper bounds on $\tan \beta_{\text{eff}}$. Since the LHC search for non-SM Higgs boson usually needs a large $\tan \beta_{\text{eff}}$ to enhance the signal rate [19–21], such upper bounds on $\tan \beta_{\text{eff}}$ may have important implication on LHC search for non-SM Higgs bosons.

We first consider the LHC search for the charged Higgs boson, which, for the charged Higgs heavier than top quark, mainly utilizes the channel $gg/gb \rightarrow t[\bar{b}]H^+$ with $H^+$ subsequently decaying to $\tau^+\nu_\tau$ [19]. In Fig. 4 we show the rate of this channel in the allowed parameter space, where the model-independent $5\sigma$ discovery sensitivity is obtained by the
ATLAS collaboration for 30 fb$^{-1}$ integrated luminosity [19]. In the calculation of the signal rate, we used the effective lagrangian method to incorporate the important SUSY corrections. Our results show that for more than 99% of the survived samples, the rate is smaller than the discovery sensitivity, which means that the LHC is unlikely to discover $H^+$. Our results also indicate that the future XENON100 limits (in case of nonobservation of DM) will further tighten the parameter space, making the discovery of $H^+$ unlikely even with higher luminosity. Note that for the charged Higgs lighter than top quark, the LHC search can instead utilize top pair production with one top decay into charged Higgs [19]. Like the case of the heavy charged Higgs boson, our results indicate that for more than 99% of the survived samples, the signal is below the 5σ discovery sensitivity obtained by the ATLAS collaboration due to $\text{Br}(t \to H^+b) < 10^{-2}$. The small branching ratio of $t \to H^+b$ arises from the fact that $\tan \beta$ is around 10 for $m_{H^+} \leq 150$ GeV (see Fig. 1) and for such a value of $\tan \beta$ there is a strong cancellation between different terms in the amplitude of this decay.

Now we turn to the LHC search for the non-SM neutral Higgs boson $H$ and $A$, for which both the ATLAS and the CMS collaborations utilize the channels $gg \to H(A)$ or $b\bar{b}H(A)$ with $H(A)$ decaying to $\tau$ leptons [19–21]. Unlike the charged Higgs boson search, for which a model-independent discovery sensitivity can be obtained, the analysis for the search of the neutral Higgs bosons is performed in certain SUSY scenarios. Here we consider the $m_{h_{\text{max}}}^\text{max}$ scenario with the following fixed parameters: $M_{\text{SUSY}} = 1$ TeV, $M_2 = 200$ GeV,
m_{\tilde{g}} = 800 \text{ GeV}, \text{ and } \bar{X}_t = A_t - \mu \cot \beta = 2 \text{ TeV. In order to show the } \mu \text{ dependence of the constraints, we choose several representative values of } \mu \text{ and scan the rest free parameters in the ranges: } 1 \leq \tan \beta \leq 80, 80 \text{ GeV} \leq m_A, m_{\tilde{t}} \leq 0.8 \text{ TeV and } 30 \text{ GeV} \leq M_1 \leq 500 \text{ GeV. In Fig. 5 we show the surviving samples on } m_A \text{ versus } \tan \beta \text{ plane together with the LHC discovery sensitivity for } 30 \text{ fb}^{-1} \text{ integrated luminosity. This sensitivity is obtained by the CMS collaboration with } H/A \rightarrow \tau^+\tau^- \rightarrow \mu + \text{jets} \text{ topology (semi-leptonic final states) [20],}
FIG. 5: Same as Fig. 3, but showing the LHC search sensitivity for the non-SM neutral Higgs bosons in the $m_h^{\text{max}}$ scenario [20]. Here the LHC sensitivity curve obtained by the CMS collaboration [20] corresponds to the $5\sigma$ discovery level, while the exclusion limit from the dark matter direct detection experiments is at 90% C.L. which is better than that obtained by the ATLAS collaboration with $H/A \rightarrow \tau^+\tau^- \rightarrow 2\ell+4\nu$ topology (full leptonic final states) [19, 21]. Note that in Fig. 5 the LHC sensitivity curve corresponds to the $5\sigma$ discovery level, while the exclusion limit from the dark matter direct detection experiments is at 90% C.L..

A few comments are due regarding the results displayed in Fig. 5:
(1) In getting these results we used the package NMSSMTools (version 2.3.1) [15] which uses micrOMEGAs (version 2.2) [22] for the calculation of the dark matter relic density. But in our calculations we extended the package by including more experimental constraints, such as the LEP search for the Higgs bosons and $B^+ \to \tau^+\nu_\tau$, so our combined constraint on the parameter space is more stringent.

(2) The CDMS-II/XENON constraints are sensitive to the value of $\mu$, i.e., as $\mu$ gets larger, the constraints become weak. The reason for this behavior is that a larger $\mu$ will result in a smaller Higgsino component in $\chi_1^0$ and hence suppress the Higgs-$\chi_1^0$-$\chi_1^0$ coupling, which will weaken the CDMS-II/XENON constraints.

(3) The LHC sensitivity for $\mu = 200$ GeV is taken directly from [20], and for other values of $\mu$ the curves are obtained by scaling the value of $\tan \beta$ so that the production rate of the Higgs bosons is same as that for $\mu = 200$ GeV. In doing this we used the package FeynHiggs2.7.1 [23] to calculate the production rate. Note that the $\mu$ parameter affects the production rate mainly by changing the $H(A)\bar{b}b$ coupling through loop correction ($\Delta_b$) which is proportional to $\mu \tan \beta/M_{SUSY}$ [12]. So the shift of the LHC sensitivity curve due to the variation of the $\mu$ value is not negligible, as shown in Fig. 5.

(4) Fig.5 shows that for $\mu = 200$ GeV no surviving samples can reach the observable level, while for larger $\mu$ values a small fraction of surviving samples can lie within the observable region. Numerically we checked that for $\mu = 400$ GeV, 700 GeV, 1 TeV, about 8%, 11% and 7% of the surviving samples lie within the observable region, respectively (for $\mu = 1$ TeV about one third of these detectable samples can even survive the future XENON100 limit). The reason for this behavior is that for a large $\mu$, although the LHC sensitivity curve is shifted upward, due to the much weakened CDMS-II/XENON constraints, some surviving samples can have quite large $\tan \beta$ values (as shown in Fig.3 and Fig.5) so that they can reach the LHC sensitivity.

(5) About the lower bound of $M_A$ as a function of $\tan \beta$, since both the $\chi N$ scattering cross section and the production rate of the Higgs boson are proportional to $\tan^2 \beta$ for large $\tan \beta$, one would naively expect the lower bound curve runs in parallel with the LHC sensitivity curve. As shown in Fig.5, this is not the case because we considered many experimental constraints and not all of them scale as $\tan^2 \beta$.  

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Conclusion: We have seen that if the MSSM is the true story, the current limits from dark matter and collider experiments already strongly constrain the parameter space, which has important implication for the LHC searches for the non-SM SUSY Higgs bosons. It turns out that in the currently allowed parameter space, the charged Higgs boson is hardly accessible at the LHC for an integrated luminosity of 30 fb$^{-1}$, while the neutral non-SM Higgs bosons ($H,A$) may be accessible in some allowed region characterized by a large $\mu$. The future XENON100 (6000 kg-days exposure) will significantly tighten the parameter space in case of nonobservation of dark matter, further shrinking the likelihood of discovering the non-SM Higgs bosons at the LHC. So the interplay of the dark matter direct detection experiments and the LHC Higgs searches will allow for a good test of the MSSM.

Finally we stress that we obtain the above conclusion by choosing a small $f_{Ts}$ ($= 0.02$). If we choose a large $f_{Ts}$, the scattering rate will be larger so that the limits from the current CDMS II/XENON will become more stringent. For example, for $f_{Ts} = 0.38$ taken in previous studies [9], we find that the current CDMS II/XENON constraints are comparable with the future XENON100 (6000 kg-days) constraints. We also checked that if we relax some assumptions in our scan, e.g., the grand unification relation for the gaugino masses, our findings about the LHC Higgs searches remain unchanged. Further, we noticed the controversy on the XENON100 detection efficiency [24]. Although our current bounds are from CDMS II plus XENON100, the CDMS II plays the dominant role. If we do not include the current XENON100 limits, our results almost remain unchanged.

Note added: After finishing our paper, we notice that the ATLAS collaboration publishes an analysis for the search sensitivity of the neutral non-SM Higgs bosons via the semi-leptonic final states [25], in which the obtained discovery sensitivity seems better than the CMS result shown in Fig.5. According to this new ATLAS result, more surviving samples in Fig.5 will reach the observable region. So our conclusion about the observability of the neutral Higgs bosons will remain unchanged.

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[1] J. Dunkley et al. [WMAP Collaboration], Astrophys. J. Suppl. 180, 306 (2009).
[2] For a review, see, e.g., H. E. Haber and G. L. Kane, Phys. Rept. 117, 75 (1985); J. F. Gunion and H. E. Haber, Nucl. Phys. B 272, 1 (1986).
[3] Z. Ahmed et al., arXiv:0912.3592.
[4] E. Aprile et al., arXiv:1005.0380.
[5] J. Hisano, K. Nakayama and M. Yamanaka, Phys. Lett. B 684, 246 (2010).
[6] M. Asano et al., arXiv:0912.5361; M. Holmes and B. D. Nelson, Phys. Rev. D 81, 055002 (2010).
[7] M. Drees and M. Nojiri, Phys. Rev. D 48, 3483 (1993).
[8] A. Djouadi and M. Drees, Phys. Lett. B 484, 183 (2000); G. Belanger et al., Comput. Phys. Commun. 180, 747 (2009).
[9] J. R. Ellis, K. A. Olive and C. Savage, Phys. Rev. D 77, 065026 (2008).
[10] 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).
[11] G. Junman, M. Kamionkowski and K. Griest, Phys. Rept. 267, 195 (1996).
[12] M. S. Carena et al., Nucl. Phys. B 577, 88 (2000).
[13] V. M. Abazov et al. [D0 Collaboration], Phys. Lett. B 660, 449 (2008).
[14] J. Cao et al., JHEP 1007, 044 (2010).
[15] U. Ellwanger, J. F. Gunion and C. Hugonie, JHEP 0502, 066 (2005).
[16] J. Cao and J. M. Yang, JHEP 0812, 006 (2008).
[17] M. Davier, et al., Eur. Phys. Jour. C 66, 1 (2010).
[18] E. Aprile and L. Baudis, PoS Identification of dark matter 2008, 018 (2008) [arXiv:0902.4253].
[19] G. Aad et al. [ATLAS Collaboration], arXiv:0901.0512.
[20] The CMS Collaboration, J. Phys. G34, 995 (2007) (also see CMS Note 2006/105).
[21] S. Horvat (on behalf of the ATLAS and CMS Collaborations), ATL-PHYS-PROC-2009-063.
[22] G. Belanger, et al., JCAP 0509, 001 (2005); Comput. Phys. Commun. 176, 367 (2007); C. Hugonie, G. Belanger, A. Pukhov, JCAP 0711, 009 (2007).
[23] S. Heinemeyer, W. Hollik and G. Weiglein, Comput. Phys. Commun. 124, 76 (2000); Eur. Phys. J. C 9, 343 (1999); G. Degrassi, et al. Eur. Phys. J. C 28, 133 (2003); M. Frank, et al. JHEP 0702, 047 (2007).

[24] J. I. Collar and D. N. McKinsey, arXiv:1005.0838; arXiv:1005.3723; The XENON100 Collaboration, arXiv:1005.2615.

[25] The ATLAS Collaboration, ATL-PHYS-PUB-2010-011.