Right-handed sneutrino dark matter in supersymmetric B - L model

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Right-handed Sneutrino Dark Matter in Supersymmetric $B - L$ Model

Shaaban Khalil
Centre for Theoretical Physics, The British University in Egypt,
El Sherouk City, Postal No. 11837, P.O. Box 43, Egypt and
Department of Mathematics, Ain Shams University, Faculty of Science, Cairo 11566, Egypt

Hiroshi Okada
Centre for Theoretical Physics, The British University in Egypt,
El Sherouk City, Postal No. 11837, P.O. Box 43, Egypt

Takashi Toma
Institute for Theoretical Physics, Kanazawa University, Kanazawa 920-1192, Japan

We show that the lightest right-handed sneutrino in TeV scale supersymmetric $B-L$ model with inverse seesaw mechanism is a viable candidate for cold dark matter. We find that it accounts for the observed dark matter relic abundance in a wide range of parameter space. The spin-independent cross section of $B - L$ right-handed sneutrino is consistent with the recent results CDMS II and XENON experiments and it is detectable in future direct detection experiments. Although the $B-L$ right-handed sneutrinos annihilate into leptons, the PAMELA results can not be explained in this model unless a huge boost factor is considered. Also the muon flux generated by $B-L$ right-handed sneutrino in the galactic center is smaller than Super-Kamiokande's upper bound.

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I. INTRODUCTION

The experimental verifications of non-vanishing neutrino masses and the alluring hints of dark matter's (DM's) existence are serious indications for new physics beyond the Standard Model (SM). Supersymmetry (SUSY) is an attractive candidate for new physics at TeV scale that provides an elegent solution for the SM gauge hierarchy problem and stabilize the SM Higgs mass at the electroweak scale. The minimal supersymmetric standard model (MSSM) is the simplest extension of the SM. In order to account for the observed neutrino masses and mixing, SM singlets (right-handed neutrinos) are usually introduced.

In the MSSM with R-parity conservation and universal soft SUSY breaking terms, the lightest neutralino, which is typically bino dominated, is an attractive candidate for cold DM. However, the current experimental constraints on SUSY particles lead to overproduction of bino relic abundance, in contradiction with the observational limits of the Wilkinson Microwave Anisotropy Probe (WMAP) [1]. In addition, the recent results of Cryogenic DM Search (CDMS II) [2] set an upper limit on the DM-nucleon elastic scattering spin independent cross section of order $3.8 \times 10^{-44}\text{cm}^2$ for DM mass of 70 GeV, which imposes stringent limits on the lightest neutralino of the MSSM, even if it consists of gaugino-Higgsino mixture. Therefore, one concludes that the DM constraints severely reduce the allowed range in the parameter space of the MSSM. It is also worth mentioning that the observed anomalies in the cosmic rays may favor the type of dark matter that annihilates into leptons not to quarks, unlike the lightest neutrlino in MSSM.

TeV scale right-handed neutrinos can be naturally implemented in supersymmeric $B - L$ extension of the SM (SUSY $B - L$), which is based on the gauge group $G_{B-L} \equiv SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ [3]. In this model, three SM singlet fermions arise quite naturally due to the $U(1)_{B-L}$ anomaly cancellation conditions. These particles are accounted for right-handed neutrinos, and hence a natural explanation for the seesaw mechanism is obtained [3–5]. This work was based on earlier papers in Ref. [6]. In this class of model, the scale of $B-L$ symmetry breaking is related to supersymmetry breaking scale [7]. Therefore, the right-handed neutrino masses are naturally of order TeV scale. This has initiated a considerable interest in analyzing the phenomenological implications of these models and their possible signatures at the LHC [8].

TeV scale SUSY $B-L$ is one of the simplest models that provides viable and testable solution to the two puzzles of the DM and the neutrino masses. However, in order to fulfill the experimental measurements for the light-neutrino masses, one of the following scenarios must be adopted: (i) Type I seesaw mechanism with very small Dirac neutrino Yukawa couplings, $Y_\nu < \mathcal{O}(10^{-7})$ [3]. (ii) Inverse seesaw mechanism with order one Yukawa couplings and small mass scale $\sim \mathcal{O}(1)$ KeV, that corresponds to scale of breaking a remnant discrete symmetry, $(-1)^{L+S}$ [10]. This work was based on earlier papers in Ref. [11]. In the first case, due to the smallness of Dirac Yukawa couplings, the right-handed neutrino sector has a very suppressed interaction with the SM particle. Therefore, the prediction of SUSY $B-L$ remains close to the MSSM ones. It turns out that the DM candidate of this model is still the lightest neutralino [12], which is a kind of a mixture of three neutral gauginos and four neutral Higgsino.

In this paper we consider the scenario where the right-handed sneutrino in SUSY $B-L$ with inverse seesaw is the lightest SUSY particle (LSP) and stable, so that it can be a cold DM candidate [13]. It is worth mentioning that in MSSM the left-handed sneutrino is the only stable weakly coupled neutral boson that can be a DM candidate [14]. However, the current limits in direct detection experiments rule out this possibility, since left-handed sneutrino has a tree level interaction with the Z gauge boson, hence its elastic scattering cross section with nucleons is quite large. In case of MSSM extension with TeV right-handed neutrino (N) superfields, the interaction of N with the SM particles can be obtained from the superpotential:

$$\mathcal{W} = \mathcal{W}_{\rm MSSM} + Y_\nu N^c L H_2. \tag{1}$$

Due to smallness of Y_ν, the annihilation cross section of the right-handed sneutrinos is very suppressed, hence its relic density is larger than the WMAP limit. Moreover, since they cannot couple to the quark sectors, the direct detection experiments such as CDMS II cannot be tested in this case.

We will show that in the SUSY $B-L$ model with $Y_\nu \sim \mathcal{O}(1)$ the dominant annihilation channel for right-handed sneutrino is given by the four point interactions, leading to h^0 and h^0. While the effective couplings of right-handed sneutrinos with quarks are obtained through the exchange of TeV scale $B-L$ gauge boson, Z_{B-L}. Thus, one finds that the WMAP results of the DM relic abundance and CDMS II/XENON for direct detection can be accommodated in SUSY $B-L$ simultaneously [15].

The paper is organized as follows. In section 2 we analyze the supersymmetric $B-L$ model with Inverse Seesaw Mechanism. We study the neutrino sector and the neutrino mass eigenstates. In section 3 we analyze the $B-L$ right-handed sneutrino mass and interactions. Section 4 is devoted for the analysis of the relic abundance of $B-L$ right-handed sneutrino. We show that the WMAP limits can be easily satisfied in a wide range of parameter space. In section 5 we discuss the direct detection rate of $B-L$ right-sneutrino DM. We show that the elastic cross section of our DM candidate with nucleon is consistent with the recent results of CDMS II/XENON experiment and it is detectable in near future experiments. In section 6 we discuss the indirect detection rate of $B-L$ right-handed sneutrino. We show that the annihilation channels of right-handed sneutrino into leptons are subdominant, therefore it cannot account for the controversial PAMELA results. We also show that the muon flux generated from the right-handed sneutrino in the galactic center is much smaller than the Super-Kamiokande's limits. Finally we give our conclusions in section 6.

II. SUPERSYMMETRIC $B-L$ MODEL WITH INVERSE SEESAW MECHANISM

The proposed TeV scale supersymmetric $B-L$ extension of the SM is based on the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$, where the $U(1)_{B-L}$ is spontaneously broken by a chiral singlet superfield χ_1 with $B-L$ charge $= +1$ and χ_2 with $B-L$ charge $= -1$. As in the conventional $B-L$ model, a gauge boson Z_{B-L} and three chiral singlet sueperfields N_i with $B-L$ charge $= -1$ are introduced for the consistency of the model. Finally, three chiral singlet superfields S_1 with $B-L$ charge $= +2$ and three chiral singlet superfields S_2 with $B-L$ charge $= -2$ are considered to implement the inverse seesaw mechanism. The superpotential of the leptonic sector of this model is given by

$$\mathcal{W} = Y_e E^c L H_1 + Y_\nu N^c L H_2 + Y_S N^c \chi_1 S_2 + \mu H_1 H_2 + \mu' \chi_1 \chi_2. \tag{2}$$

It is worth noting that the chiral singlet superfields χ_2 and N have the same $B-L$ charge. Therefore, one may impose a discrete symmetry in order to distinguish them and to prohibit other terms beyond those given in Eq. (2). In this case, the relevant soft SUSY breaking terms, assuming the usual universality assumptions, are as follows

$$\begin{aligned} -\mathcal{L}_{soft} &= \sum_\phi \widetilde{m}^2_\phi |\phi|^2 + Y^A_\nu \widetilde{N}^c \widetilde{L} H_2 + Y^A_e \widetilde{E}^c \widetilde{L} H_1 + Y^A_S \widetilde{N}^c \widetilde{S}_2 \chi_1 + B\mu H_1 H_2 + B\mu' \chi_1 \chi_2 \\ &+ \frac{1}{2} M_1 \widetilde{B}\widetilde{B} + \frac{1}{2} M_2 \widetilde{W}^a \widetilde{W}^a + \frac{1}{2} M_3 \widetilde{g}^a \widetilde{g}^a + \frac{1}{2} M_{B-L} \widetilde{Z}_{B-L} \widetilde{Z}_{B-L} + h.c, \end{aligned} \tag{3}$$

where the sum in the first term runs over $\phi = H_1, H_2, \chi_1, \chi_2, \tilde{L}, \tilde{E}^c, \tilde{N}^c, \tilde{S}_1, \tilde{S}_2$ and $Y^A_L \equiv Y_L A_L$ $(L = e, \nu, S)$ is the trilinear associated with lepton Yukawa coupling. In order to prohibit a possible large mass term $MS_1 S_2$ in the above, we assume that the particles, N^c_i, $\chi_{1,2}$, and S_2 are even under matter parity, while S_1 is an odd particle. The $B-L$ symmetry is radiatively broken by the non-vanishing vacuume expectation values (VEVs) $\langle\chi_1\rangle = v'_1$ and $\langle\chi_2\rangle = v'_2$ [7]. The tree level potential $V(\chi_1, \chi_2)$ is given by

$$V(\chi_1, \chi_2) = \frac{1}{2} g^2_{B-L} (|\chi_2|^2 - |\chi_1|^2)^2 + \mu^2_1 |\chi_1|^2 + \mu^2_2 |\chi_2|^2 - \mu^2_3 (\chi_1 \chi_2 + h.c). \tag{4}$$

At GUT scale, $\mu^2_i = m^2_0 + \mu'^2, i = 1, 2$ and $\mu^2_3 = -B\mu'$. However, they have different evolution from GUT scale to TEV scale and μ_2 becomes negative, so that $B-L$ is spontaneously broken [7]. The minimization of $V(\chi_1, \chi_2)$ leads to the following condition:

$$v'^2 = (v'^2_1 + v'^2_2) = \frac{(\mu^2_1 - \mu^2_2) - (\mu^2_1 + \mu^2_2)\cos 2\theta}{2 g^2_{B-L} \cos 2\theta}, \tag{5}$$

The angle θ is defined as $\tan\theta = v_1/v_2$. The minimization conditions also leads to

$$\sin 2\theta = \frac{2\mu^2_3}{\mu^2_1 + \mu^2_2}. \tag{6}$$

After $B-L$ breaking, the Z_{B-L} gauge boson acquires a mass [3]: $M^2_{Z_{B-L}} = 4 g^2_{B-L} v'^2$. The high energy experimental searches for an extra neutral gauge boson impose lower bounds on this mass. The stringent constraint on $U(1)_{B-L}$ obtained from LEP II result, which implies [18]

$$\frac{M_{Z_{B-L}}}{g_{B-L}} > 6 \text{ TeV}. \tag{7}$$

Now we turn to the neutrino sector and show how the observed light-neutrino masses can be obtained with $\mathcal{O}(1)$ Dirac neutrino Yukawa coupling. As can be seen from Eq. (2), after $B-L$ and EW symmetry breaking, the neutrino Yukawa interaction terms lead to the following mass terms:

$$\mathcal{L}_m^\nu = m_D \bar{\nu}_L N^c + M_N N^c S_2, \tag{8}$$

where $m_D = Y_\nu v \sin\beta$ and $M_N = Y_S v' \sin\theta$. From this Lagrangian, one can easily observe that although the lepton number is broken through the spontaneous $B-L$ symmetry breaking, a remnant symmetry: $(-1)^{L+S}$ is survived, where L is the lepton number and S is the spin. After this global symmetry is broken at much lower scale, a mass term for S_2 (and possibly for S_1 as well) is generated. Therefore, the Lagrangian of neutrino masses, in the flavor basis, is given by:

$$\mathcal{L}_m^\nu = m_D \bar{\nu}_L N^c + M_N N^c S_2 + \mu_{S_2} S_2^2 (+\mu_{S_1} S_1^2). \tag{9}$$

In the basis $\{\nu_L, N^c, S_2\}$, the 3×3 neutrino mass matrix of one generation takes the form:

$$\begin{pmatrix} 0 & m_D & 0 \\ m_D & 0 & M_N \\ 0 & M_N & \mu_{S_2} \end{pmatrix}. \tag{10}$$

The mixing matrix(O) for this mass matrix leads to the following light and heavy neutrino masses respectively in the limit of $\mu_{S_i} << m_D,\ M_N$ (where $i = 1,2$):

$$m_{\nu_\ell} = \frac{m_D^2 \mu_{S_2}}{M_N^2 + m_D^2},\ m_{\nu_{H,H'}} = \pm\sqrt{M_N^2 + m_D^2} + \frac{1}{2}\frac{M_N^2 \mu_{S_2}}{M_N^2 + m_D^2}, \tag{11}$$

where

$$O \simeq \begin{pmatrix} \frac{M_N}{\sqrt{M_N^2+m_D^2}} & \frac{1}{\sqrt{2}}\frac{m_D}{\sqrt{M_N^2+m_D^2}} + \frac{3}{4\sqrt{2}}\frac{M_N^2 m_D \mu_{S_2}}{(M_N^2+m_D^2)^2} & \frac{1}{\sqrt{2}}\frac{m_D}{\sqrt{M_N^2+m_D^2}} - \frac{3}{4\sqrt{2}}\frac{M_N^2 m_D \mu_{S_2}}{(M_N^2+m_D^2)^2} \\ \frac{M_N m_D \mu_{S_2}}{(M_N^2+m_D^2)^{3/2}} & -\frac{1}{\sqrt{2}} - \frac{1}{4\sqrt{2}}\frac{M_N^2 \mu_{S_2}}{(M_N^2+m_D^2)^{3/2}} & \frac{1}{\sqrt{2}} - \frac{1}{4\sqrt{2}}\frac{M_N^2 \mu_{S_2}}{(M_N^2+m_D^2)^{3/2}} \\ -\frac{m_D}{\sqrt{M_N^2+m_D^2}} & \frac{1}{\sqrt{2}}\frac{M_N}{\sqrt{M_N^2+m_D^2}} - \frac{1}{4\sqrt{2}}\frac{M_N(M_N^2+4m_D^2)\mu_{S_2}}{(M_N^2+m_D^2)^2} & \frac{1}{\sqrt{2}}\frac{M_N}{\sqrt{M_N^2+m_D^2}} + \frac{1}{4\sqrt{2}}\frac{M_N(M_N^2+4m_D^2)\mu_{S_2}}{(M_N^2+m_D^2)^2} \end{pmatrix}. \tag{12}$$

Thus, the light neutrino mass can be of order eV, as required by the oscillation data, for a TeV scale M_N, provided μ_{S_2} is sufficiently small, $\mu_{S_2} \ll M_N$. In this case, there is no any restriction imposed on the value of Dirac mass m_D. Therefore, the possibility of testing this type of model in LHC is quite feasible. Note that in the limit $\mu_{S_2} \to 0$ which corresponds to the unbroken $(-1)^{L+S}$ symmetry, we have massless light neutrinos. Therefore, a small non-vanishing μ_{S_2} can be considered as a slight breaking of a this global symmetry. Hence, according to 't Hooft criteria, the smallness of μ_{S_2} is natural. The possibility of generating small μ_{S_2} radiatively has been discussed in Ref. [16].

Finally, it is worth mentioning that the light neutrinos ν_l have suppressed mixing (of order $m_D\mu_{S_2}/(M_N^2 + m_D^2)$) with one type of the heavy neutrinos (say $\nu_{H'}$) and a rather small mixing (of order m_D/M_N) with the other type of heavy neutrinos (ν_H) by choosing appropriate parameters. The mixing between the heavy neutrino ν_H and ν'_H is maximal. In general, the physical neutrino states are given in terms of ν_L, N^c, and S_2 as follows:

$$\nu_l \simeq \frac{M_N}{\sqrt{M_N^2+m_D^2}}\nu_L + \frac{M_N m_D \mu_{S_2}}{(M_N^2+m_D^2)^{3/2}} N^c - \frac{m_D}{\sqrt{M_N^2+m_D^2}} S_2 \simeq \nu_L + a_1 N^c - a_2 S_2 \tag{13}$$

$$\nu_H \simeq -\frac{1}{\sqrt{2}}\frac{m_D}{\sqrt{M_N^2+m_D^2}}\nu_L + \frac{1}{\sqrt{2}} N^c - \frac{1}{\sqrt{2}}\frac{M_N}{\sqrt{M_N^2+m_D^2}} S_2 \simeq \alpha(-a_2\ \nu_L +\ N^c -\ S_2) \tag{14}$$

$$\nu_{H'} \simeq \frac{1}{\sqrt{2}}\frac{m_D}{\sqrt{M_N^2+m_D^2}}\nu_L + \frac{1}{\sqrt{2}} N^c + \frac{1}{\sqrt{2}}\frac{M_N}{\sqrt{M_N^2+m_D^2}} S_2 \simeq \alpha(a_2\nu_L +\ N^c +\ S_2). \tag{15}$$

For $m_D \simeq 100$ GeV, $M_N \simeq 1$ TeV and $\mu_{S_2} \simeq 1$ KeV, one finds that $a_1 \sim \mathcal{O}(10^{-10})$, $a_2 \sim \mathcal{O}(0.1)$, $a_3 \sim \mathcal{O}(0.07)$ and $\alpha \sim \sin\pi/4$. In this respect, the gauge eigenstates for neutrinos can be expressed in terms of the mass eigenstates as follow:

$$\begin{pmatrix} \nu_L \\ N^c \\ S_2 \end{pmatrix} = O \begin{pmatrix} \nu_l \\ \nu_H \\ \nu_{H'} \end{pmatrix} \simeq \begin{pmatrix} 1 & \alpha a_2 & \alpha a_2 \\ a_1 & -\alpha & \alpha \\ -a_2 & \alpha & \alpha \end{pmatrix} \begin{pmatrix} \nu_l \\ \nu_H \\ \nu_{H'} \end{pmatrix}. \tag{16}$$

III. $B-L$ RIGHT-HANDED SNEUTRINO

In our model, the sneutrino mass matrix of one generation is given by 8×8 matrix, which can be decomposed into the following two mass matrices: (i) 6×6 mass matrix in basis of $(\tilde{\nu}_L, \tilde{\nu}_L^\dagger, \tilde{N}, \tilde{N}^\dagger, \tilde{S}_2, \tilde{S}_2^\dagger)^T$. (ii) 2×2 mass matrix in basis of $(\tilde{S}_1, \tilde{S}_1^\dagger)^T$. The $\tilde{S}_1$'s are decoupled and have now interactions with the SM particles. Therefore, one can neglect it and focus on the 6×6 sneutrino mass matrix. In the flavor basis; $\tilde{\nu} \equiv (\tilde{\nu}_L, \tilde{\nu}_L^\dagger, \tilde{N}, \tilde{N}^\dagger, \tilde{S}_2, \tilde{S}_2^\dagger)^T$, the sneutrino mass matrix $\tilde{M}_{\tilde{\nu}}^2$ is given as

$$M_{\tilde{\nu}}^2 = \begin{pmatrix} M^2_{\tilde{\nu}_L^\dagger \tilde{\nu}_L} & 0 & (M^2_{\tilde{N}^\dagger \tilde{\nu}_L})^\dagger & 0 & (M^2_{\tilde{S}_2^\dagger \tilde{\nu}_L})^\dagger & 0 \\ 0 & (M^2_{\tilde{\nu}_L^\dagger \tilde{\nu}_L})^T & 0 & (M^2_{\tilde{N}^\dagger \tilde{\nu}_L})^T & 0 & (M^2_{\tilde{S}_2^\dagger \tilde{\nu}_L})^T \\ M^2_{\tilde{N}^\dagger \tilde{\nu}_L} & 0 & M^2_{\tilde{N}^\dagger \tilde{N}} & 0 & (M^2_{\tilde{S}_2^\dagger \tilde{N}})^\dagger & (M^2_{\tilde{S}_2 \tilde{N}})^\dagger \\ 0 & (M^2_{\tilde{N}^\dagger \tilde{\nu}_L})^* & 0 & (M^2_{\tilde{N}^\dagger \tilde{N}})^T & (M^2_{\tilde{S}_2 \tilde{N}})^T & (M^2_{\tilde{S}_2^\dagger \tilde{N}})^T \\ M^2_{\tilde{S}_2^\dagger \tilde{\nu}_L} & 0 & M^2_{\tilde{S}_2^\dagger \tilde{N}} & (M^2_{\tilde{S}_2 \tilde{N}})^* & M^2_{\tilde{S}_2^\dagger \tilde{S}_2} & (M^2_{\tilde{S}_2 \tilde{S}_2})^\dagger \\ 0 & (M^2_{\tilde{S}_2^\dagger \tilde{\nu}_L})^* & M^2_{\tilde{S}_2 \tilde{N}} & (M^2_{\tilde{S}_2^\dagger \tilde{N}})^* & M^2_{\tilde{S}_2 \tilde{S}_2} & (M^2_{\tilde{S}_2^\dagger \tilde{S}_2})^T \end{pmatrix}, \tag{17}$$

where

$$M^2_{\tilde{\nu}_L^\dagger \tilde{\nu}_L} = \tilde{m}^2_{\nu_L} + v^2 \cos^2\beta Y^{e\dagger} Y^e + M_D^\dagger M_D + \frac{m_Z^2}{2}\cos 2\beta - \frac{M^2_{Z_{B-L}}}{4}(1-\cot^2\theta), \tag{18}$$

$$M^2_{\tilde{N}^\dagger \tilde{N}} = \tilde{m}^2_N + v'^2 \sin^2\theta Y^S Y^{S\dagger} + M_D M_D^\dagger - \frac{M^2_{Z_{B-L}}}{4}(1-\cot^2\theta), \tag{19}$$

$$M^2_{\tilde{S}_2^\dagger \tilde{S}_2} = \tilde{m}^2_{S_2} + |\mu_{S_2}|^2 + M_N^\dagger M_N - \frac{M^2_{Z_{B-L}}}{2}(1-\cot^2\theta), \tag{20}$$

$$M^2_{\tilde{N}^\dagger \tilde{\nu}_L} = \mu^* v \cos\beta Y^\nu + v \sin\theta Y^\nu_A, \quad M^2_{\tilde{S}_2^\dagger \tilde{\nu}_L} = v v' \sin\theta \sin\beta Y^{S\dagger} Y^\nu, \tag{21}$$

$$M^2_{\tilde{S}_2^\dagger \tilde{N}} = \mu' v' \cos\theta Y^{S\dagger} + v' \sin\theta Y_A^{\nu\dagger}, \quad M^2_{\tilde{S}_2 \tilde{S}_2} = B_2' \mu_{S_2}, \quad M^2_{\tilde{S}_2 \tilde{N}} = \mu_{S_2} v' \sin\theta Y_S^\dagger. \tag{22}$$

In the case of existing the mixing of $\tilde{S}_2\tilde{S}_2$ or $\tilde{S}_2\tilde{N}$, the complex scalar DM splits in two real scalar and the lighter one is DM. If the mass split is small as well as the momentum of DM, inelastic scattering can be considered.

The mass matrix is diagonalized by unitary matrix Γ as

$$\Gamma^\dagger M_{\tilde{\nu}}^2 \Gamma = \mathrm{diag}(m^2_{\tilde{\nu}_1^m}, m^2_{\tilde{\nu}_2^m}, m^2_{\tilde{\nu}_3^m}, m^2_{\tilde{\nu}_4^m}, m^2_{\tilde{\nu}_5^m}, m^2_{\tilde{\nu}_6^m}). \tag{23}$$

Thus, the mass eigenstates $\tilde{\nu}^m$ are defined as $\tilde{\nu}_i = \Gamma_{ij}\tilde{\nu}_j^m$. In general, the lightest sneutrino can be written as a linear combination of the sneutrino mass eigenstate, as follows:

$$\tilde{\nu}_1^m = \Gamma^\dagger_{11}\tilde{\nu}_L + \Gamma^\dagger_{12}\tilde{\nu}_L^\dagger + \Gamma^\dagger_{13}\tilde{N} + \Gamma^\dagger_{14}\tilde{N}^\dagger + \Gamma^\dagger_{15}\tilde{S}_2 + \Gamma^\dagger_{16}\tilde{S}_2^\dagger. \tag{24}$$

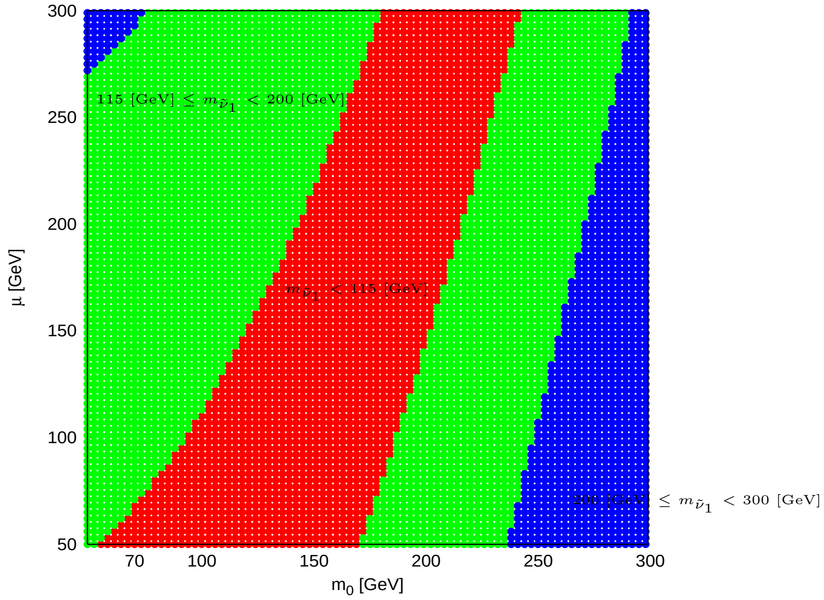

FIG. 1: A contour plot for the mass of $B-L$ right-handed sneutrino in terms of the most relevant parameters: soft mass term m_0 and μ term and for $\theta \simeq \beta \simeq \pi/4$ and $Y_\nu \simeq 1$. The red region represents $m_{\tilde{\nu}_1} < 115$ GeV, which is kinematically excluded because DM mass is lighter than the SM Higgs mass. The green region represents 115 GeV $\leq m_{\tilde{\nu}_1} < 200$ GeV. The blue region represents 200 GeV $\leq m_{\tilde{\nu}_1} < 300$ GeV.

However if one considers large $\tan\beta$ and small $\tan\theta$ limits, one finds that the lightest sneutrino is mainly obtained from the $(\tilde{N}, \tilde{S}_2)$ sector. Therefore, it can be expressed as

$$\tilde{\nu}_1 \simeq \Gamma^\dagger_{13}\tilde{N} + \Gamma^\dagger_{14}\tilde{N}^\dagger + \Gamma^\dagger_{15}\tilde{S}_2 + \Gamma^\dagger_{16}\tilde{S}^\dagger_2. \tag{25}$$

The mass of this particle, which we call "$B-L$ right-handed sneutrino" depends on the universal soft scalar mass m_0 and on the parameters μ and μ'. In Fig. 1, we display the mass range of $B-L$ right-handed sneutrino as function of the most relevant parameters m_0 and μ. If this particle is the lightest SUSY particle, then it is stable and can be considered as an interesting candidate for DM.

Now we consider the relevant interactions of the $B-L$ right-handed sneutrino. From the superpotential $\mathcal{W}$ in Eq. (2) one gets the following interacting Lagrangian of $\tilde{N}$ in the flavor basis:

$$\begin{aligned}\mathcal{L}^W_{int} &= Y_{\nu ij}\tilde{N}^\dagger_i\left[(\tilde{H}^0_2)^c P_L \nu_{Lj} - (\tilde{H}^+_2)^c P_L \ell^-_{Lj}\right] + Y_{Sij}\tilde{N}^\dagger_i(\tilde{S}_{2j})^c P_L\widetilde{\chi}_1 \\ &+ Y^A_{\nu ij}\tilde{N}^\dagger_i\left(\tilde{\nu}_{Lj}H^0_2 - \tilde{\ell}^-_{Lj}H^+_2\right) + Y^A_{Sij}\tilde{N}^\dagger_i\tilde{S}_{2j}\chi_1 + h.c.,\end{aligned} \tag{26}$$

Also from the F-term contributions to the scalar potential one finds the following interaction terms:

$$\begin{aligned}\mathcal{L}^F_{int} &= -|H^-_1|^2\tilde{\nu}^\dagger_{Li}(Y^\dagger_e Y_e)_{ij}\tilde{N}_j - (|H^+_2|^2 + |H^0_2|^2)\tilde{N}^\dagger_i(Y^\dagger_\nu Y_\nu)_{ij}\tilde{N}_j \\ &- (H^{0\dagger}_1 H^+_2 + H^{-\dagger}_1 H^0_2)\tilde{N}^\dagger_i(Y^\nu Y^{e\dagger})_{ij}\tilde{\ell}_{Rj} - (\tilde{N}^\dagger Y_\nu\tilde{\nu}_L)(\tilde{\nu}^\dagger_L Y^\dagger_\nu\tilde{N}) - \mu^* H^{0\dagger}_1(\tilde{N}^\dagger Y_\nu\tilde{\nu}_L) \\ &- (\tilde{N}^\dagger Y_\nu\tilde{\ell}_L)(\tilde{\ell}^\dagger_L Y^\dagger_\nu\tilde{N}) - \mu^* H^{-\dagger}_1(\tilde{N}^\dagger Y_\nu\tilde{\ell}_L) - (\tilde{S}^\dagger_2 Y^\dagger_S\tilde{N})(\tilde{N}^\dagger Y_S\tilde{S}_2) - \mu^*\chi^\dagger_2(\tilde{N}^\dagger Y_S\tilde{S}_2) \\ &- |\chi_1|^2\tilde{N}^\dagger_i(Y_S Y^\dagger_S)\tilde{N}_j - \mu^*_{S_2}\chi_1\tilde{N}^\dagger_i Y_{Sij}\tilde{S}^\dagger_{2j} + h.c..\end{aligned} \tag{27}$$

Next the interactions of $\tilde{N}$ with the gauge fields lead to the following Lagrangian:

$$\mathcal{L}^G_{int} = g^2_{B-L}Z^2_{B-L}\tilde{N}^\dagger_i\tilde{N}_i - ig_{B-L}Z^\mu_{B-L}\left(\tilde{N}_i\partial_\mu\tilde{N}^\dagger_i - \tilde{N}^\dagger_i\partial_\mu\tilde{N}_i\right) - i\sqrt{2}g_{B-L}\tilde{N}_i N^c_i\tilde{Z}_{B-L} + h.c.. \tag{28}$$

Finally, the D-term implies that

$$\begin{aligned}\mathcal{L}^{D}_{int} &= g^2_{B-L}(\tilde{N}^\dagger\tilde{N}) \\ &\times\left(-|\tilde{\nu}_L|^2-|\tilde{N}|^2+\frac{|\tilde{u}_L|^2}{3}+\frac{|\tilde{d}_L|^2}{3}+\frac{|\tilde{u}_R|^2}{3}+\frac{|\tilde{d}_R|^2}{3}-|\tilde{\ell}_L|^2-|\tilde{\ell}_R|^2+2|\tilde{S}_1|^2-2|\tilde{S}_2|^2+|\tilde{\chi}_1|^2-|\tilde{\chi}_2|^2\right). \end{aligned} \tag{29}$$

In our analysis for $B-L$ right-handed sneutrino ($\tilde{\nu}_1$) DM, the following assumptions are considered for simplification: (i) Each sector consists of one generation only. (ii) $m_{SM} < m_{DM} < m_{SUSY}$, $m_{h^\pm}$, $m_{\chi_{1,2}}$, $M_{Z_{B-L}}$, $m_{S_{1,2}}$, where m_{SM} is standard model particles including of the lightest neutral Higgs boson (h_0), m_{SUSY} is the supersymmetric particles, and m_{DM} is the mass of the lightest SUSY sneutrinos. (iii) In chargino sector, $h^\pm$ mass is approximately given by μ, in the limit of μ, $M_2 >> m_W$, where m_W is the SM charged weak gauge boson mass. Under these assumptions, the relevant interacting Lagrangian of $\tilde{\nu}_1$ consists of

$$\mathcal{L}_{int} = \mathcal{L}^W_{int} + \mathcal{L}^F_{int} + \mathcal{L}^G_{int}. \tag{30}$$

Note that the D-term is kinematically irrelevant, while the W, F, G interactions take the following simplified form:

$$\mathcal{L}^W_{int} = Y'_\nu \Gamma_{41}\tilde{\nu}_1^\dagger\left[\overline{\tilde{h}^{0c}}P_L\nu_L - \overline{\tilde{h}^{+c}}P_L\ell^-_L\right] + h.c., \tag{31}$$

$$\mathcal{L}^F_{int} = -(Y^\dagger_\nu Y_\nu)\Gamma_{41}\Gamma_{31}|h^0|^2\tilde{\nu}_1^\dagger\tilde{\nu}_1, \tag{32}$$

$$\mathcal{L}^G_{int} = -ig_{B-L}Z^\mu_{B-L}\Gamma_{41}\Gamma_{31}\left(\tilde{\nu}_1\partial_\mu\tilde{\nu}_1^\dagger - \tilde{\nu}_1^\dagger\partial_\mu\tilde{\nu}_1\right), \tag{33}$$

where Y'_ν is considered to be included in the mixing of Higgsino and Chargino. Hereafter we use it as Y_ν. Eq. (31) might be relevant to the indirect detection, which will be discussed in the section VI. Eq. (32) is more relevant to the WMAP experiment, which will be discussed in the next section. Eq. (33) is applied to analyze the direct detection as CDMS II/XENON, which will be also discussed in the section V.

IV. RELIC ABUNDANCE OF $B-L$ RIGHT-HANDED SNEUTRINO

In this section, we compute the relic abundance of $B-L$ right-handed sneutrino DM. We consider the standard computation of the cosmological abundance, where $\tilde{\nu}_1$ is assumed to be in thermal equilibrium with the SM particles in the early universe and decoupled when it was non-relativistic. Therefore, the $\tilde{\nu}_1$ density can be obtained by solving the Boltzmann equation:

$$\frac{dn_{\tilde{\nu}_1}}{dt} + 3Hn_{\tilde{\nu}_1} = - < \sigma^{ann}_{\tilde{\nu}_1} v > [(n_{\tilde{\nu}_1})^2 - (n^{eq.}_{\tilde{\nu}_1})^2], \tag{34}$$

where $n_{\tilde{\nu}_1}$ is $\tilde{\nu}_1$ number density with $\rho_{\tilde{\nu}_1} = m_{\tilde{\nu}_1}n_{\tilde{\nu}_1}$. One usually defines $\Omega_{\tilde{\nu}_1} = \rho_{\tilde{\nu}_1}/\rho_c$, where ρ_c is the critical mass density. In addition, $< \sigma^{ann}_{\tilde{\nu}_1} v >$ is the thermal averaged of the total cross section for $\tilde{\nu}_1$ annihilation into SM lighter particles times the DM relative velocity v. For non-relativistic $\tilde{\nu}_1$, the thermal averaged annihilation cross section, $< \sigma^{ann}_{\tilde{\nu}_1} v >$, can be approximated as follows [22]:

$$< \sigma^{ann}_{\tilde{\nu}_1} v > \simeq a_{\tilde{\nu}_1} + b_{\tilde{\nu}_1} v^2, \tag{35}$$

where $a_{\tilde{\nu}_1}$ $b_{\tilde{\nu}_1}$ are the coefficients coming from s-wave and p-wave of $\tilde{\nu}_1\tilde{\nu}_1$ annihilation, respectively.

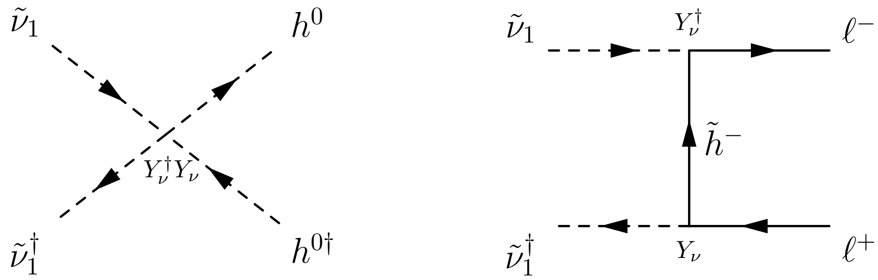

FIG. 2: Possible annihilation channels of $\tilde{\nu}_1$. The second diagram gives a sub-dominant contribution, however it may be relevant for indirect detection processes.

From Eqs.(31-33), one finds that the dominant annihilation channels of $\tilde{\nu}_1$ are given in Fig. 2. It turns out the annihilation of $B-L$ right-handed sneutrino into SM-like neutral Higgs, through the four point interaction vertex, gives the dominant contribution. The tree level annihilation channel $\tilde{\nu}_1\tilde{\nu}_1 \to \ell^+\ell^-$ is suppressed by the mass of the chargino exchanged particle. This channel may be relevant for the indirect detection processes which will be discussed later. Our computation for the annihilation cross section leads to the following $a_{\tilde{\nu}_1}$ $b_{\tilde{\nu}_1}$:

$$a_{\tilde{\nu}_1} = \frac{\beta'_{h^0}}{32\pi m^2_{\tilde{\nu}_1}} |Y_\nu \Gamma_{31}\Gamma_{41}|^4, \tag{36}$$

$$b_{\tilde{\nu}_1} = \frac{\beta'_{h^0}(x^2_{h^0}-1)}{128\pi m^2_{\tilde{\nu}_1}} |Y_\nu \Gamma_{31}\Gamma_{41}|^4, \tag{37}$$

where

$$z_a = \frac{m_a}{m_{\tilde{\nu}_1}}, \quad \beta'^2_a = 1 - z_a^2,\ x_a^2 = \frac{z_a^2}{2(1-z_a^2)}, \tag{38}$$

here we define that a is a final-state particle. From Eq.(34) one finds that the relic abundance $\Omega_{\tilde{\nu}_1}h^2$ is given by [23]

$$\Omega_{\tilde{\nu}_1}h^2 \simeq \frac{8.76\times 10^{-11}\mathrm{GeV}^{-2}}{g_*^{1/2}(T_F)(a_{\tilde{\nu}_1}/x_F + 3b_{\tilde{\nu}_1}/x_F^2)}, \tag{39}$$

where

$$x_F = \ln \frac{0.0955\ m_{\rm pl}\ m_{\tilde{\nu}_1}(a_{\tilde{\nu}_1} + 6b_{\tilde{\nu}_1}/x_F)}{(g_*^{1/2}(T_F)x_F)^{1/2}}. \tag{40}$$

Here $m_{\rm pl}$ is the Planck mass (1.22×10^{19} GeV) and $g_*(T_F)$ enumerates the degrees of freedom of relativistic particles at the freeze out temperature T_F, which can be fixed as $g_*^{1/2}(T_F) = 10$. From the above expressions, one notes that $\tilde{\nu}_1$ relic abundance depends only on the right-sneutrino mass and the annihilation cross section coefficients $a_{\tilde{\nu}_1}$ and $b_{\tilde{\nu}_1}$.

Given $a_{\tilde{\nu}_1}$ $b_{\tilde{\nu}_1}$ we can determine the freeze out temperature T_F, below which the $B-L$ right-handed sneutrino annihilation rate is smaller than the expansion rate of the universe and then computing the relic density $\Omega_{\tilde{\nu}_1}h^2$. In our numerical computation we assume a universal soft mass (m_0) and fix the other parameter as follows:

$$m_{h^0} = 115\ \mathrm{GeV},\ M_{Z_{B-L}} = 6\ \mathrm{TeV},\quad M_N = 1\ \mathrm{TeV},\ \mu_{S_{1,2}} = 1\ \mathrm{KeV},\ v = 175\ \mathrm{GeV},$$
$$m_D = 100\ \mathrm{GeV},\ B'_{1,2} = v' = \mu^{(')} = 1\ \mathrm{TeV},\quad Y_S = 0.1,\ Y^e = 0.05,\ Y_\nu^A = Y_S^A = 0.1\ \mathrm{TeV}, \tag{41}$$

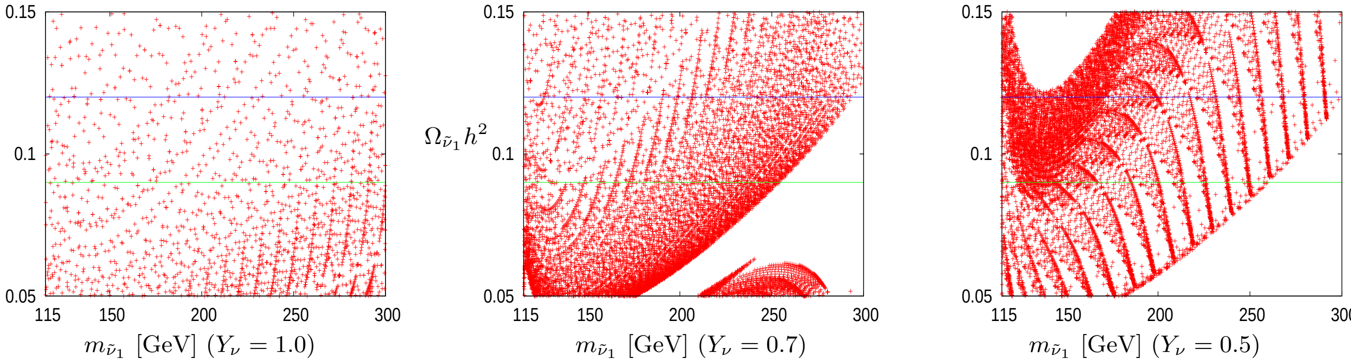

FIG. 3: Relic density $\Omega_{\tilde{\nu}_1} h^2$ as function of the $B-L$ right-handed sneutrino mass $m_{\tilde{\nu}_1}$ for three values of Dirac neutrino Yukawa coupling. The region between two lines is allowed by the experiment of WMAP [1].

In addition we analyze the relic density of $\tilde{\nu}_1$ in the following regions of the parameter μ, β, and θ:

$$0 < \theta \;,\; \beta < \pi, \qquad 50 \text{ GeV} < \mu \;,\; m_0 < 300 \text{ GeV}. \tag{42}$$

In Fig. 3, we present the values of relic abundance of $B-L$ right-handed sneutrino $\Omega_{\tilde{\nu}_1} h^2$ as a function of $m_{\tilde{\nu}_1}$ for the following values of Dirac neutrino Yukawa coupling: $Y_\nu = 0.5, 0.7, 1.0$. We require $B-L$ right-handed sneutrino relic density to be $0.09 < \Omega_{\tilde{\nu}_1} h^2 < 0.12$ in order to be consistent with WMAP results at 3σ [24]. As can be seen from this figure, smaller values of Dirac neutrino Yukawa coupling are favored and lead to more allowed points that satisfy the WMAP observational limits of DM relic density. Also we find that the smaller Y_ν is considered, the smaller DM mass one obtains.

V. DIRECT DETECTION

In this section we discuss the possibility to detect our $B-L$ right-handed sneutrino in direct detection experiments such as CDMS (II) [2] and XENON 100 experiment [25]. The general form of the elastic scattering cross section between DM $\tilde{\nu}_1$ and nuclei N is given by [22, 44]

$$\sigma^{\text{vec}}_{\tilde{\nu}_1 - N} = \frac{m_r^2}{16\pi} \left| b_N \right|^2 . \tag{43}$$

The reduced mass m_r is defined as

$$m_r = \left(\frac{1}{m_{\tilde{\nu}_1}} + \frac{1}{M} \right)^{-1} , \tag{44}$$

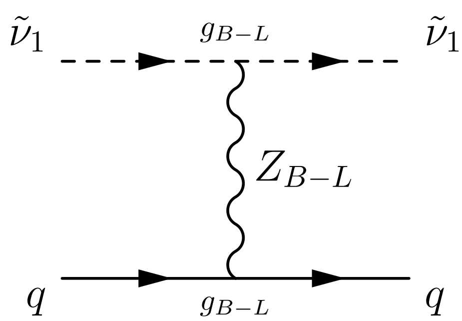

FIG. 4: Our dominant diagram for the direct detection

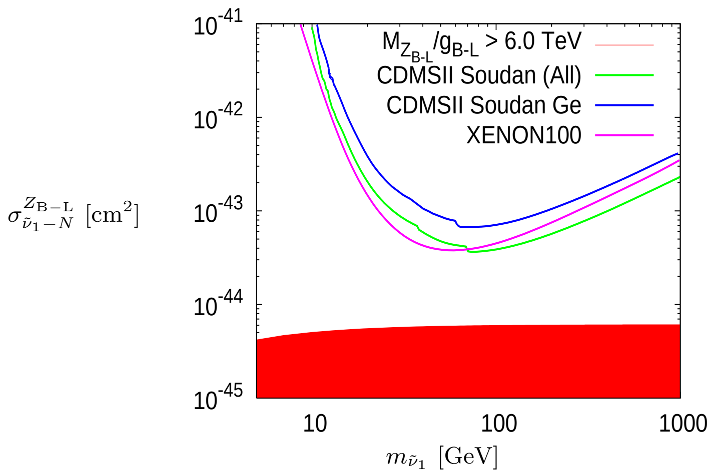

FIG. 5: Elastic scattering cross section between a nuclei and the $B-L$ right-handed sneutrino DM as function of the DM mass for $g_{B-L}=1$.

where M is the nuclei mass. The coefficient b_N is given by

$$b_N = (A-Z)b_n + Zb_p, \quad b_p = 2b_u + b_d, \quad b_n = b_u + 2b_d, \tag{45}$$

Here A and Z are the mass number and the atomic number, respectively. The effective Lagrangian parameters b_u and b_d are defined as

$$\mathcal{L}_{\rm eff} = b_q X^\mu \bar{q}\gamma^\mu q, \quad q=(u\ ,\ d). \tag{46}$$

Here X^μ is a general form of the vector current. In case of fermionic DM X^μ is given by $X^\mu \simeq \bar{f}\gamma^\mu f$. While for bosonic DM, it is defined as $X^\mu \simeq ib^\dagger\partial^\mu b - ib\partial^\mu b^\dagger$.

In our $B-L$ case, the elastic scattering cross section of the right-handed sneutrino with a given nuclei has a spin-independent contribution arising from Z_{B-L} gauge boson exchange diagrams, as can been seen in Fig. 4. The interactions between the right-handed sneutrino $\tilde{\nu}_1$ and Z_{B-L} boson in $B-L$ model are

$$\mathcal{L} \ \supset \ -ig_{B-L}Z^\mu_{B-L}\Gamma_{41}\Gamma_{31}\left(\tilde{\nu}_1\partial_\mu\tilde{\nu}_1^\dagger - \tilde{\nu}_1^\dagger\partial_\mu\tilde{\nu}_1\right) - \frac{1}{3}g_{B-L}Z^\mu_{B-L}\bar{u}\gamma_\mu u - \frac{1}{3}g_{B-L}Z^\mu_{B-L}\bar{d}\gamma_\mu d. \tag{47}$$

Writing down the effective interaction, one obtains

$$i\mathcal{L}_{\rm eff} \supset -\frac{|\Gamma_{41}\Gamma_{31}|^2}{M^2_{Z_{B-L}}}\frac{g^2_{B-L}}{3}\left(\tilde{\nu}_1\partial_\mu\tilde{\nu}_1^\dagger - \tilde{\nu}_1^\dagger\partial_\mu\tilde{\nu}_1\right)\bar{u}\gamma^\mu u - \frac{|\Gamma_{41}\Gamma_{31}|^2}{M^2_{Z_{B-L}}}\frac{g^2_{B-L}}{3}\left(\tilde{\nu}_1\partial_\mu\tilde{\nu}_1^\dagger - \tilde{\nu}_1^\dagger\partial_\mu\tilde{\nu}_1\right)\bar{d}\gamma^\mu d. \tag{48}$$

Assuming $|\Gamma_{41}\Gamma_{31}|^2 \simeq 1$ for simplicity, one finds

$$b_p = b_n = i\frac{g^2_{B-L}}{M^2_{Z_{B-L}}}. \tag{49}$$

Therefore, b_N is given by

$$b_N = iA\frac{g^2_{B-L}}{M^2_{Z_{B-L}}}. \tag{50}$$

Thus, the elastic scattering cross section of $B-L$ right-handed sneutrino is given by

$$\sigma^{Z_{B-L}}_{\tilde{\nu}_1-N} = \frac{m_r^2}{16\pi}|A|^2\frac{g^4_{B-L}}{M^4_{Z_{B-L}}}. \tag{51}$$

In Fig. 5 we depict the relation between the cross section and $B-L$ right-sneutrino mass. As can been seen from this figure, the following upper bound on $\sigma^{Z_{B-L}}_{\tilde{\nu}_1-N}$ can be obtained:

$$\sigma^{Z_{B-L}}_{\tilde{\nu}_1-N} \leq 6.2\times 10^{-45}\ \mathrm{cm}^2. \qquad (52)$$

It is also remarkable that the elastic cross section is quite insensitive to the $B-L$ right-handed sneutrino mass $m_{\tilde{\nu}_1}$. However, one observes that a light sneutrino $\sim \mathcal{O}(100)$ GeV is more favored by direct detection experimental results. The current limits from CDMS II and XENON experiments indicate to a lower-bound of order 3.7×10^{-44} cm^2. This suggests that our $B-L$ right-handed sneutrino DM is expected to be detected in the direct detection experiments in near future.

Before concluding this section, it is worth mentioning that the XENON 100 [26] experiment has recently presented new limits on the WIMP-nucleon cross section for inelastic DM. These limits are due to a DM run with 100.9 live-days of data, taken from January to June 2010. It was shown that $\sigma_{\chi-N} < 10^{-41}$ cm^2 can be extracted for $m_\chi > 100$ GeV. Where χ is a generic dark matter. The bound rules out the explanation of controversial DAMA/LIBRA modulation results, as being due to inelastic DM. The bound obtained on $\sigma_{\chi-N}$ from the inelastic DM analysis of XENON 100 should be considered carefully, since a minimum velocity for DM to scatter in a detector is introduced, hence a large amount of fiducial is needed. Nevertheless, as can be seen from Eq. (52), the cross section of our $B-L$ sneutrino is well below the XENON 100 bound.

VI. INDIRECT DETECTION

A. PAMELA and Fermi-LAT experiments

As advocated in the introduction, the indirect searches for DM by the Space Observatory PAMELA [34] and Fermi-LAT [35], indicate that the DM may contribute to the positron flux by direct annihilation in $\ell^+\ell^-$. This is one of the main feature of our DM candidate $B-L$ right-handed sneutrino, therefore, it important to investigate the possibility that it accounts for these results. PAMELA collaboration reported excess flux between 8 and 80 GeV, with no excess in the corresponding anti-proton flux. Also ATIC and Fermi-LAT balloon experiments have shown excess electron and positron flux at energies around $10-1000$ GeV. While there are plausible astrophysical explanations for these excesses, such as local pulsars and supernovas remnants, they could also result from DM annihilation. Note that in order to explain the Fermi-LAT experiment by DM annihilation, the DM mass must be of order $\mathcal{O}$(TeV). As shown in the previous section, such heavy DM mass is not favored by of direct detections. Therefore, in our analysis, we assume that the Fermi-LAT experiment may be saturated by considering an astrophysical background and we will focus on PAMELA measurement.

It is known that if the DM annihilation is to explain the observed anomalous flux, a large annihilation cross section, $\langle\sigma v\rangle \simeq 10^{-24}$ cm^{3}s^{-1}, is required to fit the excess flux, which is incompatible with straightforward estimates of the relic DM abundance in conventional cosmological models. Otherwise, a huge, unexplained, boost factor must be introduced. In our SUSY $B-L$ model, the $B-L$ right-handed sneutrino annihilates into $\ell^+\ell^-$ channels, as shown in the second diagram of Fig. 2. However, as discussed in the previous section that these channels give sub-dominate contribution to the annihilation process. Therefore, the corresponding annihilation cross section is $< 10^{-27}$ cm^{3}s^{-1}. In this case, one requires a huge boost factor $\mathcal{O}(10^{5-6})$ at least in order to account for PAMELA results. In general, it is known that there are two mechanisms to enhance the cross section and may justify this large boost factor. The first is Breit-Wigner mechanism [36] and the second is the Sommer-feld [37] enhancement. Breit-Wigner mechanism can not be implemented in our model, since there is no any diagrams with s-channel. On the other hand, the Sommer-feld enhancement requires higher DM mass compared to the mediated particles in order to obtain enough large boost factor $> \mathcal{O}(10^{5-6})$. This assumption we have already avoided in order to not spoil the direct detection. As a result, it is difficult for our $B-L$ right-handed neutrino to explain the controversial results of PAMELA experiment.

B. Muon flux measurement from Super-Kamiokande

The high energy neutrinos induced by DM annihilations in the earth, the sun, and the galactic center is an important signal for the indirect detection of DM. Such energetic neutrinos induce upward through-going muons from charged current interactions provide the most effective signatures in Super-Kamiokande. The neutrino-induced muon flux is evaluated from the neutrino flux [40, 41] as

$$F^{(\mathrm{ann})}_{\mu^+\mu^-} \simeq 5.9\times 10^{-15}\ \mathrm{cm}^{-2}\mathrm{s}^{-1} \times \sum_F S_F \left(\frac{\langle\sigma v\rangle_F}{10^{-23}\ \mathrm{cm}^3\mathrm{s}^{-1}}\right)\left(\frac{\langle J_2\rangle_\Omega \Delta\Omega}{10}\right). \tag{53}$$

where we fix the typical values $\langle J_2\rangle_\Omega\Delta\Omega \sim 10$ for $\psi_{\max} = 5^\circ$ in case of the Navarro-Frenk-White (NFW) halo density profile [42], and F collectively denotes the primary annihilation mode (e.g., $\tau^+\tau^-$, etc.). Notice that model dependence comes from $\langle\sigma v\rangle_F$, which will be shown later. S_F is defined as

$$S_F = \sum_{\nu_i}\int_{E_{\min}}^{E_{\mathrm{in}}} \frac{dN_F^{(\nu_i)}}{dE} P_{\nu_i\nu_\mu}\left(\frac{E}{E_{\mathrm{in}}}\right)^2 dE, \tag{54}$$

where $E_{\mathrm{in}} = m_{\tilde{\nu}_1}$, and $E_{\min}$ is the threshold energy above which the muons can be detected. $P_{\nu_i\nu_\mu}$ denotes the probability that the ν_i at the production is observed as ν_μ at the Earth due to the effect of neutrino oscillation. Regardless to the complicated expression of S_F, it is found as a fixed value depending on each of the F particle, which is, e.g., 0.2 for μ pair, $0.14-0.18$ for τ pair, 0.78 for ν_τ pair, etc [38]. The limits from Super-Kamiokande are given in the ref. [43].

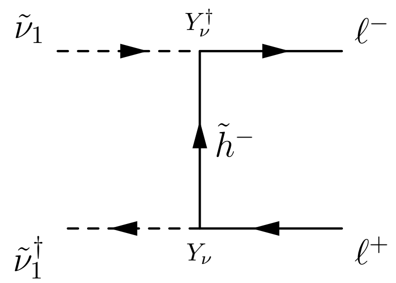

FIG. 6: Fenymen diagram of $\tilde{\nu}_1\tilde{\nu}_1 \to \ell^-\ell^+$. Note that, in general, there is an another contribution of $\tilde{\nu}_1, \tilde{\nu}_1 \to \tilde{h}^0 \to \nu, \nu$ process. However, due to the masslless limit of ν, the cross section of this channel vanishes.

In SUSY $B-L$ model, the relevant Lagrangian in given by

$$\mathcal{L}^W_{int} = Y_\nu \Gamma_{41}\tilde{\nu}_1^\dagger \left[\overline{\tilde{h}^{0c}} P_L \nu_L - \overline{\tilde{h}^{+c}} P_L \ell^-_L\right] + h.c. \tag{55}$$

One can show that in the limit $v \to 0$, the thermally averaged cross section is given by

$$\begin{aligned} \langle \sigma v \rangle_F |_{v\to 0} &\simeq \sum_{F=\tau,\nu_\tau} \frac{3\beta'_F}{16\pi m^2_{\tilde{\nu}_1}} |Y_\nu \Gamma_{31}\Gamma_{41}|^4 z_F^2 (2-z_F^2) \times \left(r^2_{F,\tilde{h}^+} + r^2_{F,\tilde{h}^-} + r_{F,\tilde{h}^-} r_{F,\tilde{h}^+} + r^2_{F,\tilde{h}^0} \right) \\ &\simeq \sum_{F=\tau} \frac{9\beta'_F}{16\pi m^2_{\tilde{\nu}_1}} |Y_\nu \Gamma_{31}\Gamma_{41}|^4 z_F^2 (2-z_F^2) \times r^2_{F,\tilde{h}^+}. \end{aligned} \tag{56}$$

Note that the term proportional to $r^2_{F,\tilde{h}^0}$ vanishes in the massless limit of ν. The parameters z_F and β'_F are as defined in Eq. (38). The other parameters are defined as follows:

$$w_\alpha = \frac{m_\alpha}{m_{\tilde{\nu}_1}}, \qquad r_{F\alpha} = (1 - z_F^2 + w_\alpha^2)^{-1}. \tag{57}$$

Here F refers to the final-state particle and α denotes the mediated-particle. In our analysis for the muon flux, we consider the same set of inputs that we have used in the previous sections that leads to relic abundance within the WMAP limits: $0.09 < \Omega h^2 < 0.12$. Moreover, since we are considering the effect one generation, we assume τ final state only. Using S_F that corresponds to τ final state, one finds $\langle \sigma v \rangle$

$$\begin{aligned} F^{(\rm ann)}_{\mu^+\mu^-} &\simeq 5.9 \times 10^{-15}\ {\rm cm^{-2} s^{-1}} \times \sum_{F=\tau^\pm} S_F \left(\frac{\langle \sigma v \rangle_F}{8.56 \times 10^{-7}} \right) \left(\frac{\langle J_2 \rangle_\Omega \Delta\Omega}{10} \right) \\ &\simeq 6.9 \times 10^{-9}\ {\rm cm^{-2} s^{-1}} \times \sum_{F=\tau^\pm} S_F \left(\frac{9\beta'_F}{16\pi m^2_{\tilde{\nu}_1}} |Y_\nu \Gamma_{31}\Gamma_{41}|^4 z_F^2 (2-z_F^2) \times r^2_{F,\tilde{h}^+} \right). \end{aligned} \tag{58}$$

As we fixed the cone-half angle from the galactic center: $\psi_{\rm max} = 5^\circ$ which is maximum in the case of NFW profile, one finds that the Super-Kamiokande limit should be less than $5 \times 10^{-15} {\rm cm^{-2} sec^{-1}}$ (See figures of Ref. [38]).

In Fig. 7, we plot the muon flux induced from the $B - L$ annihilation in the galactic center. As can be seen from this figure, the result of the induced muon flux is $\lesssim 10^{-20}\ {\rm cm^{-2} sec^{-1}}$, which is few order of magnitudes smaller than the upper bound of Super-Kamiokande. This result is expected since the thermally averaged cross cross is much smaller than the typical value required by indirect detections.

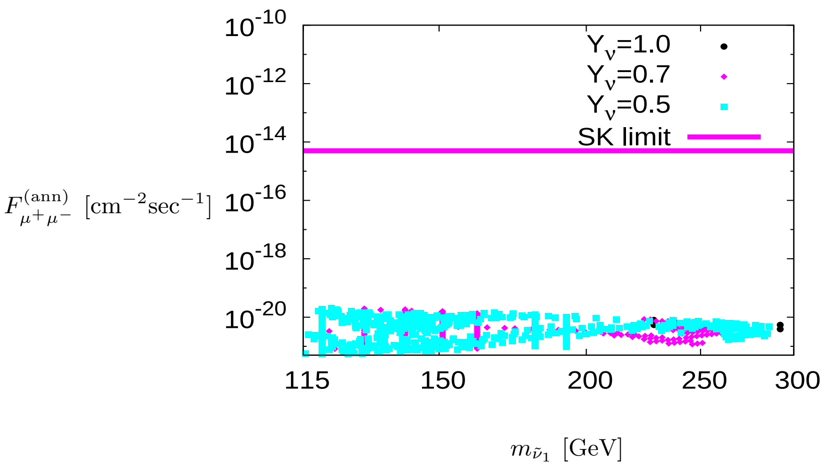

FIG. 7: Muon flux induced from the $B - L$ annihilation in the galactic center. The horizontal line refers to the Super-Kamiokande upper bound which is given by $\sim 5 \times 10^{-15} {\rm cm^{-2} sec^{-1}}$.

VII. CONCLUSIONS

In this paper we considered the supersymmetric $B-L$ model with inverse seesaw mechanism. We demonstrated that the lightest right-handed sneutrino in this model can be stable and a viable candidate for cold dark matter. We studied the relic abundance of the $B-L$ right-handed sneutrino and showed that the WMAP result, $\Omega h^2 \simeq 0.11$, can be satisfied in a wide range of the parameter space. We emphasized that the dominate annihilation channel of $B-L$ right-handed sneutrino is given by $\tilde{\nu}_1\tilde{\nu}_1 \to h^0 h^0$, where h^0 is the SM-like Higgs. We also studied the direct detection rate of $B-L$ right-handed sneutrino. We found that its elastic cross section is consistent with the upper bounds of current experiments, such as CDMS II and XENON. Our result of $B-L$ right-handed sneutrino direct detection is promising and indicates that it can be detectable in near future experiment.

In addition, we have analyzed the indirect detection rate of $B-L$ right-handed sneutrino. In particular, we focused on the observation of the Space Observatory PAMELA for positron flux and also on the neutrino-induced upward through-going muons in the Super-Kamiokande detector. We showed that although the $B-L$ right-handed sneutrino annihilates at tree level into leptons, the corresponding cross section is much smaller than the required one for accommodating PAMELA results. Also the neutrino flux induced by the $B-L$ right-handed sneutrino annihilations in the galactic center is much smaller than the Super-Kamiokande's limits.

Finally it is worth mentioning that our $B-L$ right-handed sneutrino may be produced at the Large Hadron Collider (LHC) through the channel is $q\bar{q} \to Z_{B-L} \to \tilde{\nu}_1^\dagger \tilde{\nu}_1$. However, the amplitude of this channel vanishes identically due to the fact that left and right quarks or leptons has the same $B-L$ quantum numbers. However, slepton/left-handed sneutrino may decay to right-handed sneutrino, which escapes the detector and gives the missing energy signal similar to other examples of cold DM [46].

Acknowledgments

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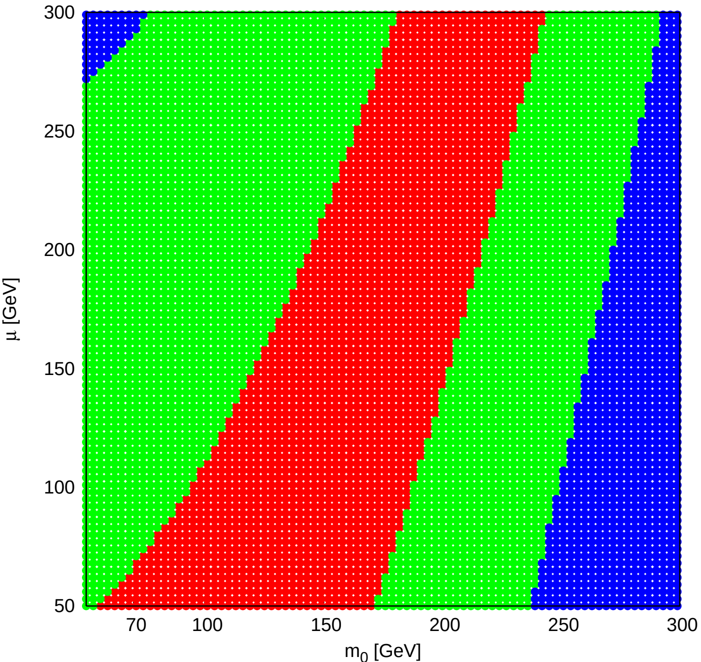
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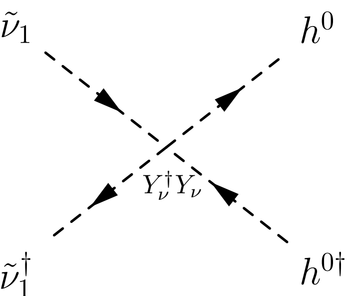
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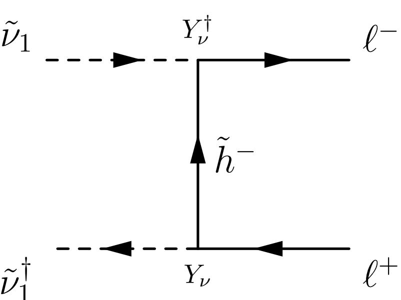
h^0
$Y_\nu^\dagger Y_\nu$
$\tilde{\nu}_1^\dagger$
$h^{0\dagger}$



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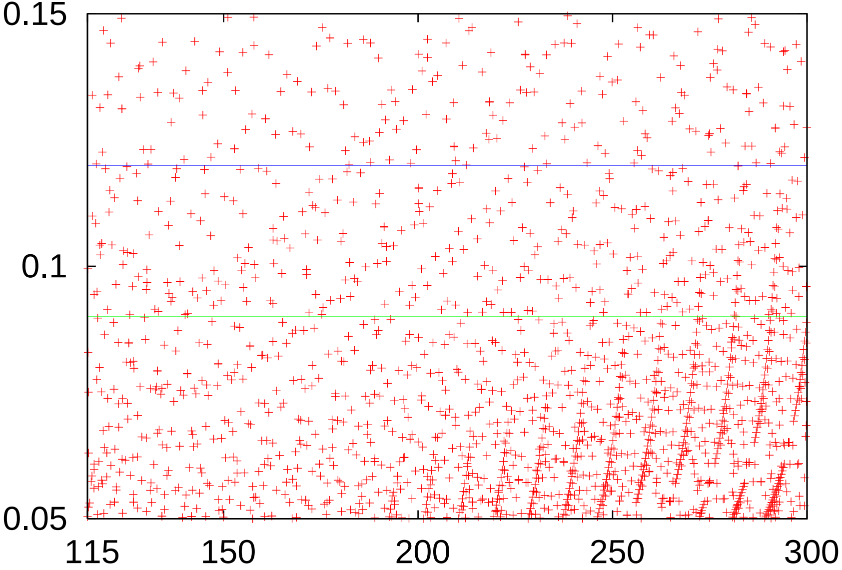
Y_ν
ℓ^+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0.1
0.05
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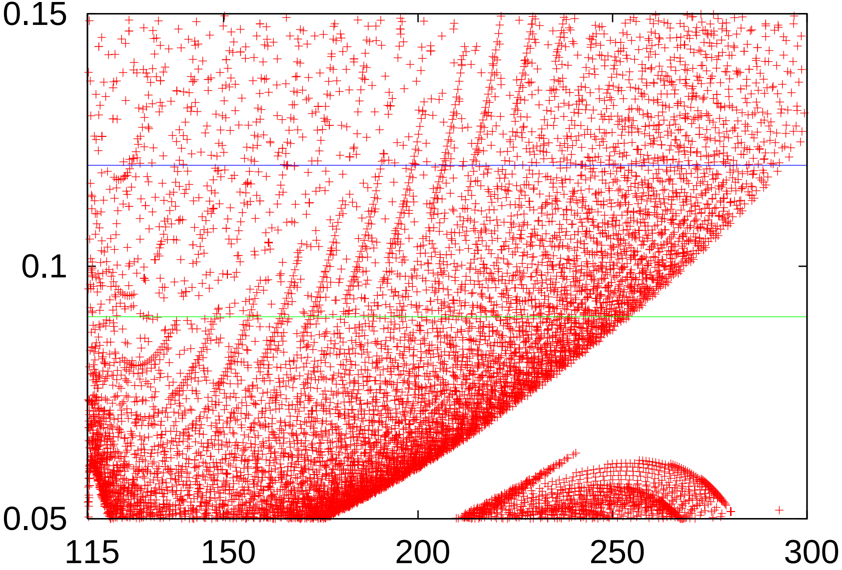

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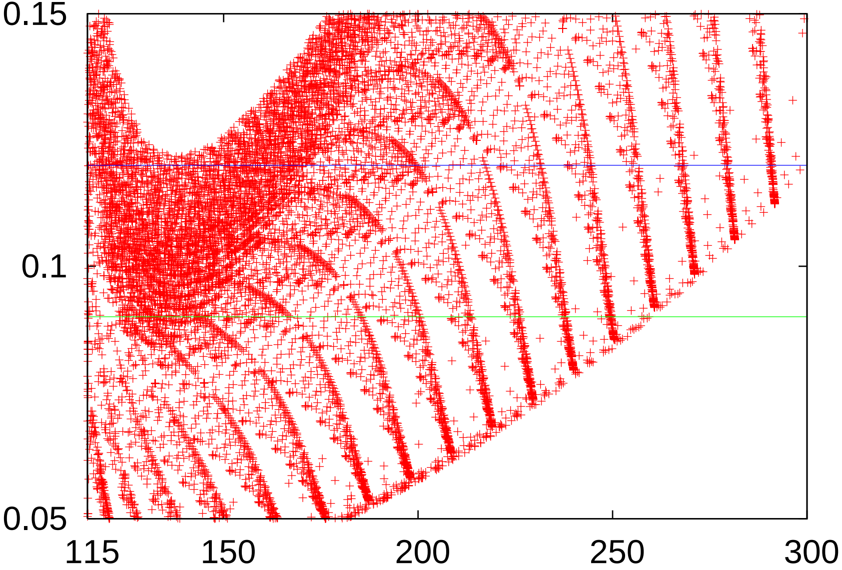
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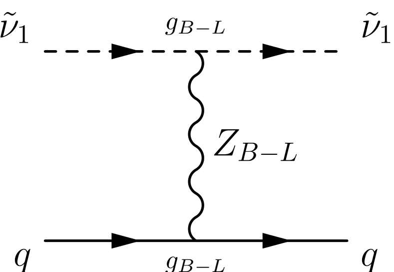
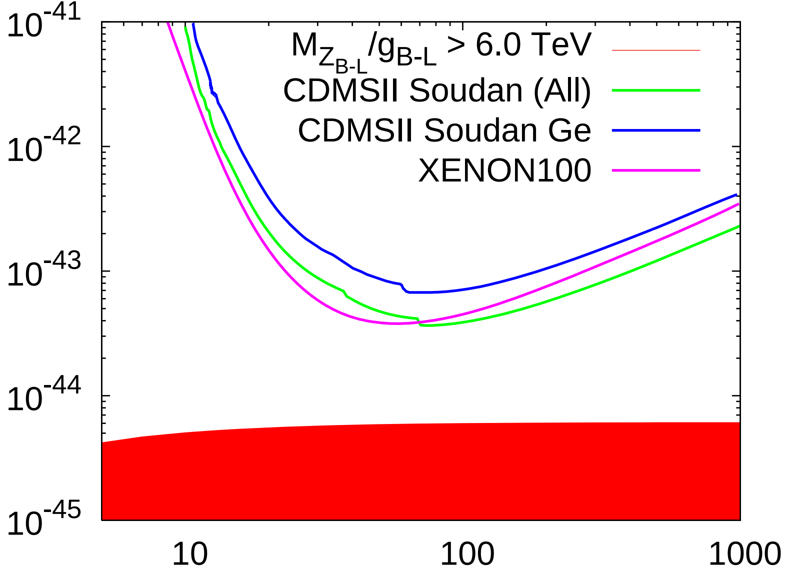
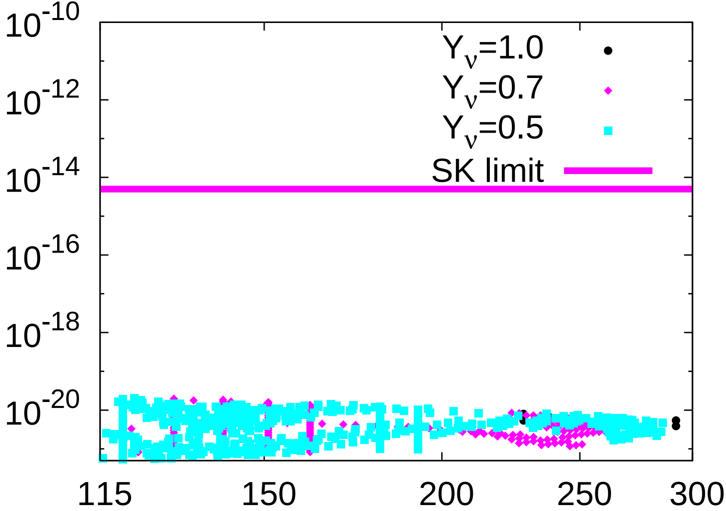


$\tilde{\nu}_1$
g_{B-L}
$\tilde{\nu}_1$
Z_{B-L}
q
g_{B-L}
q



$M_{Z_{B-L}}/g_{B-L}$ > 6.0 TeV
CDMSII Soudan (All)
CDMSII Soudan Ge
XENON100
10^{-41}
10^{-42}
10^{-43}
10^{-44}
10^{-45}
10
100
1000



10^{-10}
10^{-12}
10^{-14}
10^{-16}
10^{-18}
10^{-20}
Y_ν=1.0
Y_ν=0.7
Y_ν=0.5
SK limit
115
150
200
250
300