Supersymmetric minimal B-L model at the TeV scale with right-handed Majorana neutrino dark matter

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Abstract

We propose a supersymmetric extension of the minimal $B-L$ model where we consider a new Z_2-parity under which one right-handed neutrino is assigned odd parity. When the Majorana Yukawa coupling of a Z_2-even right-handed neutrino is large, radiative corrections will drive the mass squared of the corresponding right-handed sneutrino to negative values, breaking the $B-L$ gauge symmetry at the TeV scale in a natural way. Additionally, R-parity is broken and thus the conventional supersymmetric dark matter candidate, the neutralino, is no longer viable. Thanks to the Z_2-parity, the Z_2-odd right-handed neutrino remains a stable dark matter candidate even in the presence of R-parity violation. We demonstrate that the dark matter relic abundance with an enhanced annihilation cross section by the $B-L$ gauge boson (Z') resonance is in accord with the current observations. Therefore, it follows that the mass of this dark matter particle is close to half of the Z' boson mass. If the Z' boson is discovered at the Large Hadron Collider, it will give rise to novel probes of dark matter: The observed Z' boson mass will delineate a narrow range of allowed dark matter mass. If the Z' boson decays to a pair of dark matter particles, a precise measurement of the invisible decay width can reveal the existence of the dark matter particle.

1 Introduction

The minimal supersymmetric (SUSY) extension of the Standard Model (MSSM) is one of the prime candidates for physics beyond the Standard Model (SM), which naturally solves problems in the SM, in particular, the gauge hierarchy problem. In addition, a candidate for the cold dark matter, which is missing in the SM, is also naturally incorporated in the MSSM. Searching for SUSY is one of the major occupations of the Large Hadron Collider (LHC) resources. The LHC is operating at unprecedented luminosities, and are collecting data very rapidly. Discovery of physics beyond the Standard Model in the near future is highly likely, and anticipated.

The MSSM can solve the gauge hierarchy problem and the dark matter problem, which it is able to achieve merely by virtue of it being supersymmetric. However it is clear that the SUSY extension is not enough to solve the aforementioned problems in addition to explaining neutrino phenomena. This is because the solar and atmospheric neutrino oscillation have established non-zero neutrino masses and mixings between different neutrino flavors [1]. Unlike the quark sector, the scale of neutrino masses is very small and the different flavors are largely mixed. We have no choice but to extend the MSSM in order to incorporate such neutrino masses and flavor mixings. The seesaw extension [2] has gained much attention since it not only accounts for the neutrino mass but also explains the smallness of the mass in a natural way. Corresponding to the seesaw scale (the typical scale of right-handed neutrinos) being, for example, from 1 TeV to 10^{14} GeV, the scale of the neutrino Dirac mass varies from 1 MeV (the electron mass scale) to 100 GeV (the top quark mass scale).

The $B-L$ (baryon number minus lepton number) is an anomaly-free global symmetry in the SM and it can be easily gauged. The minimal $B-L$ model is the simplest gauged $B-L$ extension of the SM [3], where right-handed neutrinos of three generations and a Higgs field with two units of the $B-L$ charge are introduced. The existence of the three right-handed neutrinos is crucial in canceling the gauge and gravitational anomalies. In this model, the mass of right-handed neutrinos arises associated with the $B-L$ gauge symmetry breaking.

Although the scale of the $B-L$ gauge symmetry breaking is arbitrary as long as phenomenological constraints are satisfied, it is interesting to consider it at the TeV scale [4]. For example, it has been recently pointed out [5] that when the classical conformal invariance is imposed on the minimal $B-L$ model, the symmetry breaking scale appears to be the TeV scale naturally. If this is the case, all new particles in the model, the Z' gauge boson, the $B-L$ Higgs boson and the right-handed neutrinos appear at the TeV scale, which can be discovered at the LHC [6]. The minimal $B-L$ model also has interesting cosmological prospects such as dark matter physics [7] and baryogenesis [8].

In this paper, we investigate supersymmetric extension of the minimal $B-L$ model. It has been pointed out [9] that in this model the $B-L$ symmetry is radiatively broken by the interplay between large Majorana Yukawa couplings of right-handed neutrinos and the soft SUSY breaking masses, a mechanism that is analogous to radiative electroweak symmetry breaking in the MSSM. The mechanism naturally places the $B-L$ symmetry breaking scale at the TeV scale.

Despite this remarkable feature of the SUSY minimal $B-L$ model, a more thorough analysis [10] indicated that most of the $B-L$ symmetry breaking parameter space is occupied by non-zero vacuum expectation values (VEVs) from right-handed sneutrinos. Therefore, the most likely scenario in the SUSY minimal $B-L$ model with the radiative $B-L$ symmetry breaking, is that R-parity is violated in the vacuum. This means that the lightest superpartner (LSP) neutralino, which is the conventional dark matter candidate in SUSY models, becomes unstable and no longer remains a viable dark matter candidate. As discussed in [11], even though R-parity is broken, an unstable gravitino if it is the LSP has a lifetime longer than the age of the universe and can still be the dark matter candidate. Although this is an interesting possibility, a mechanism of SUSY breaking mediations providing us with the LSP gravitino is limited, and we do not consider the LSP gravitino in this paper.

Recently, a cogent framework for dark matter was elucidated in the context of the (non-SUSY) minimal $B-L$ model [7], where a new Z_2-parity was introduced and one right-handed neutrino was assigned odd Z_2-parity while the other fields were assigned even Z_2. Calculation of the relic abundance of the Z_2-odd right-handed neutrino showed that it could account for the observed relic abundance, and therefore the dark matter in our universe. We mention this to emphasize that we are not introducing any new particle in the current model.

In this paper, we apply the same idea to the SUSY generalization of the minimal $B-L$ model with the radiative $B-L$ symmetry breaking, and investigate the resulting phenomenology. What we discovered is that the $B-L$ gauge symmetry and R-parity are both broken at the TeV scale by the non-zero VEV of a Z_2-even right-handed sneutrino, for suitable regions of parameter space. Even in the presence of R-parity violation, the Z_2-parity is still exact and the stability of the Z_2-odd right-handed neutrino is guaranteed. Therefore, the Z_2-odd right-handed neutrino appears to be a natural, stable dark matter candidate. We calculated the relic abundance of the Z_2-odd right-handed neutrino and found that the resultant relic abundance was in agreement with observations.

This paper is organized as follows. In the next section, we define the SUSY minimal $B-L$ model with Z_2-parity and introduce superpotential and soft SUSY breaking terms relevant for our discussion. In Sec. 3, we perform a numerical analysis of the renormalization group equation (RGE) evolution of the soft SUSY breaking masses of the right-handed sneutrinos and $B-L$

Higgs fields and show that the $B-L$ gauge symmetry is radiatively broken at the TeV scale. It will be shown that one Z_2-even right-handed sneutrino develops a VEV and hence R-parity is also radiatively broken. In Sec. 4, we calculate the relic abundance of the right-handed neutrino and identify the parameter region consistent with the observed dark matter relic abundance. We also discuss phenomenological constraints of the model in Sec. 5. The last section is devoted for conclusions and discussions.

2 Supersymmetric minimal $B-L$ model with Z_2-parity

The minimal $B-L$ extended SM is based on the gauge group $\mathrm{SU}(3)_c\times\mathrm{SU}(2)_L\times\mathrm{U}(1)_Y\times\mathrm{U}(1)_{B-L}$ with three right-handed neutrinos and one Higgs scalar field with $B-L$ charge 2 but which is a singlet under the SM gauge group[1]. As far as the motivation to introduce three generations of right-handed neutrinos (N_i^c) is concerned, the introduction of the three generations of right-handed neutrinos is in no way ad-hoc; On the contrary, once we gauge $B-L$, their introduction is forced upon us by the requirement of the gauge and gravitational anomaly cancellations. The SM singlet scalar works to break the $\mathrm{U}(1)_{B-L}$ gauge symmetry by its VEV and at the same time, generates Majorana masses for right-handed neutrinos which then participate in the seesaw mechanism.

It is easy to supersymmetrize this model and the particle contents are listed in Table 1[2]. The gauge invariant superpotential relevant for our discussion is given by

$$W_{BL} = \sum_{i=2}^{3}\sum_{j=1}^{3} y_D^{ij} N_i^c L_j H_u + \sum_{k=1}^{3} y_k \Phi N_k^c N_k^c - \mu_\Phi \bar{\Phi}\Phi, \tag{1}$$

where the first term is the neutrino Dirac Yukawa coupling, the second term is the Majorana Yukawa coupling for the right-handed neutrinos, and a SUSY mass term for the SM singlet Higgs fields is given in the third term. Without loss of generality, we have worked in the basis where the Majorana Yukawa coupling matrix is real and diagonal. Note that Dirac Yukawa couplings between N_1^c and L_j are forbidden by the Z_2-parity, so that the lightest component field in N_1^c is stable, as long as the Z_2-parity is exact.

As we will discuss in the next section, the $B-L$ gauge symmetry is radiatively broken at the TeV scale, and the right-handed neutrinos obtain TeV-scale Majorana masses. The seesaw mechanism[3] sets the mass scale of light neutrinos at $m_\nu = \mathcal{O}(y_D^2 v_u^2/M_R)$, where v_u is the VEV

[1] Recently, a global $B-L$ extended model with a dark matter candidate has been proposed [12], which has some interesting differences to the gauged model.

[2] It is possible to construct a phenomenologically viable SUSY $B-L$ model without Φ and $\bar{\Phi}$ [13].

[3] As we will see in the next section, R-parity is also radiatively broken. In this case, the right-handed neutrinos mix with the $B-L$ gaugino and fermionic components of $\bar{\Phi}$ and Φ, and the seesaw formula is quite involved.

chiral superfield	$SU(3)_c$	$SU(2)_L$	$U(1)_Y$	$U(1)_{B-L}$	R-parity	Z_2
Q^i	$\mathbf{3}$	$\mathbf{2}$	$+1/6$	$+1/3$	$-$	$+$
U_i^c	$\mathbf{3}^*$	$\mathbf{1}$	$-2/3$	$-1/3$	$-$	$+$
D_i^c	$\mathbf{3}^*$	$\mathbf{1}$	$+1/3$	$-1/3$	$-$	$+$
L_i	$\mathbf{1}$	$\mathbf{2}$	$-1/2$	-1	$-$	$+$
N_1^c	$\mathbf{1}$	$\mathbf{1}$	0	$+1$	$-$	$-$
$N_{2,3}^c$	$\mathbf{1}$	$\mathbf{1}$	0	$+1$	$-$	$+$
E_i^c	$\mathbf{1}$	$\mathbf{1}$	-1	$+1$	$-$	$+$
H_u	$\mathbf{1}$	$\mathbf{2}$	$+1/2$	0	$+$	$+$
H_d	$\mathbf{1}$	$\mathbf{2}$	$-1/2$	0	$+$	$+$
Φ	$\mathbf{1}$	$\mathbf{1}$	0	-2	$+$	$+$
$\bar{\Phi}$	$\mathbf{1}$	$\mathbf{1}$	0	$+2$	$+$	$+$

Table 1: Particle contents: In addition to the MSSM particles, three right-handed neutrino superfields ($N^c_{1,2,3}$) and two Higgs superfields ($\bar{\Phi}$ and Φ) are introduced. The Z_2-parity for N^c_1 is assigned to be odd. $i = 1, 2, 3$ is the generation index.

of the up-type Higgs doublet in the MSSM, and $M_R = \mathcal{O}(1 \text{ TeV})$ is the mass scale of the right-handed neutrinos. It is natural to assume that the mass of the heaviest light neutrino is $m_\nu \sim \sqrt{\Delta m^2_{23}} \sim 0.05$ eV with $\Delta m^2_{23} \simeq 2.43 \times 10^{-3}$ eV^2 being the atmospheric neutrino oscillation data [1]. Thus, we estimate $y_D \sim 10^{-6}$, and point out that such a small neutrino Dirac Yukawa coupling is negligible in the analysis of RGEs.

Next, we introduce soft SUSY breaking terms for the fields in the $B-L$ sector:

$$\begin{aligned} \mathcal{L}_{\rm soft} &= -\left(\frac{1}{2}M_{BL}\lambda_{BL}\lambda_{BL} + h.c.\right) - \left(\sum_{k=1}^{3} m^2_{\tilde{N}^c_k}|\tilde{N}^c_k|^2 + m^2_\Phi|\Phi|^2 + m^2_{\bar{\Phi}}|\bar{\Phi}|^2\right) \\ &+ \left(B_\Phi\bar{\Phi}\Phi + \sum_{k=1}^{3} A_k\Phi\tilde{N}^c_k\tilde{N}^c_k + h.c.\right). \end{aligned} \quad (2)$$

Here we have omitted terms relevant to the neutrino Dirac Yukawa couplings since they are very small, i.e. $O(10^{-6})$ or smaller. For simplicity, in this analysis we consider the same setup as the constrained MSSM and assume the universal soft SUSY breaking parameters, $m^2_{\tilde{N}^c_k} = m^2_\Phi = m^2_{\bar{\Phi}} = m^2_0$ and $A_k = A_0$, at the grand unification scale[4], $M_U = 2 \times 10^{16}$ GeV.

Before closing this section, we comment on the uniqueness of the Z_2-parity assignment in the phenomenological point of view. One may find the Z_2-parity assignment ad-hoc, but we cannot assign an odd-parity for any MSSM particles because the parity forbids the Dirac Yukawa couplings which is necessary to reproduce the observed fermion masses and quark flavor

[4]However, we do not necessarily assume grand unification behind our model. In fact, it is very non-trivial to unify the Z_2-odd right-handed neutrino with Z_2-even fields.

mixings. As we will see in the next section, the scalars Φ and $\bar{\Phi}$ develop non-zero VEVs to break the $B-L$ gauge symmetry, and these fields should be Z_2-parity even in order to keep the parity unbroken. Hence, we can assign Z_2-odd parity only for right-handed neutrinos. Considering the fact that we need at least two right-handed neutrinos to reproduce the observed neutrino oscillation data, two right-handed neutrinos should be parity even and be involved in the seesaw mechanism. As a result, we have assigned Z_2-parity odd for only one right-handed neutrino as in Table 1.

3 Radiative $B-L$ symmetry breaking and R-parity

In the non-SUSY minimal $B-L$ model, the $B-L$ symmetry breaking scale is determined by parameters in the Higgs potential which can in general be at any scale as long as the experimental constraints are satisfied. The LEP experiment has set the lower bound on the $B-L$ symmetry breaking scale as $m_{Z'}/g_{BL} \geq 6-7$ TeV [14]. Recent LHC results for Z' boson search with 1.1 fb^{-1} [15] excluded the $B-L$ Z' gauge boson mass $m_{Z'} \lesssim 1.5$ TeV [16] when the $B-L$ coupling is not too small. We see that the LEP bound is more severe than the LHC bound for $m_{Z'} \gtrsim 1.5$ TeV. The SUSY extension of the model, however, offers a very interesting possibility for constraining the $B-L$ symmetry breaking scale, as pointed out in [9].

It is well-known that the electroweak symmetry breaking in the MSSM is triggered by radiative corrections to the up-type Higgs doublet mass squared via the large top Yukawa coupling [17]. Directly analogous to this situation, the $B-L$ symmetry breaking occurs through radiative corrections with a large Majorana Yukawa coupling.

We consider the following RGEs for soft SUSY breaking terms in the $B-L$ sector [10, 18] :

$$
\begin{aligned}
16\pi^2\mu\frac{dM_{BL}}{d\mu} &= 48g_{BL}^2 M_{BL}, \\
16\pi^2\mu\frac{dm^2_{\tilde{N}^c_i}}{d\mu} &= 8y_i^2 m_\Phi^2 + 16y_i^2 m^2_{\tilde{N}^c_i} + 8A_i^2 - 8g_{BL}^2 M_{BL}^2, \\
16\pi^2\mu\frac{dm^2_{\Phi}}{d\mu} &= 4\left(\sum_{i=1}^{3} y_i^2\right) m_\Phi^2 + 8\sum_{i=1}^{3} y_i^2 m^2_{\tilde{N}^c_i} + 4\sum_{i=1}^{3} A_i^2 - 32g_{BL}^2 M_{BL}^2, \\
16\pi^2\mu\frac{dm^2_{\bar{\Phi}}}{d\mu} &= -32g_{BL}^2 M_{BL}^2, \\
16\pi^2\mu\frac{dA_i}{d\mu} &= \left(30y_i^2 + 2\sum_{j\neq i} y_j^2 - 12g_{BL}^2\right) A_i + 4y_i\left(\sum_{j\neq i} y_j A_j - 6g_{BL}^2 M_{BL}\right), \qquad (3)
\end{aligned}
$$

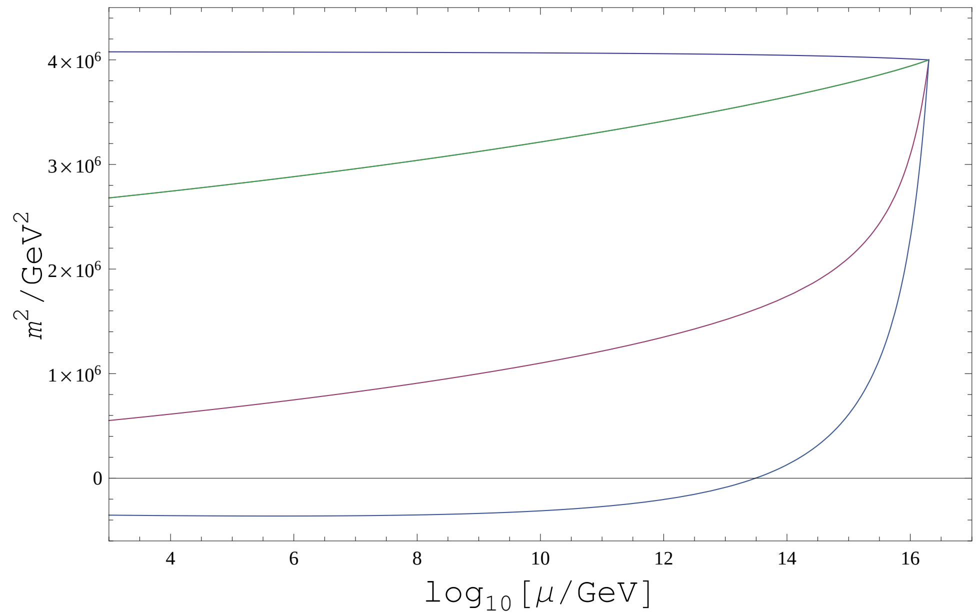

Figure 1: The RGE running of the soft SUSY breaking masses, $m^2_{\bar{\Phi}}$, $m^2_{\tilde{N}^c_1} = m^2_{\tilde{N}^c_2}$, m^2_{Φ} and $m^2_{\tilde{N}^c_3}$ from above.

where RGEs for the gauge and Yukawa couplings are given by

$$\begin{aligned} 16\pi^2\mu\frac{dg_{BL}}{d\mu} &= 24g^3_{BL}, \\ 16\pi^2\mu\frac{dy_i}{d\mu} &= y_i\left(10y_i^2 + 2\sum_{j\neq i} y_j^2 - 12g^2_{BL}\right). \end{aligned} \tag{4}$$

To illustrate the radiative B-L symmetry breaking, we solve these equations from $M_U = 2\times 10^{16}$ GeV to low energy, choosing the following boundary conditions.

$$\begin{aligned} & g_{BL} = 0.532, \;\; y_1 = y_2 = 0.4, \; y_3 = 2.5, \\ & M_{BL} = 500 \text{ GeV}, \;\; m_{\tilde{N}^c_i} = m_\Phi = m_{\bar{\Phi}} = 2 \text{ TeV}, \;\; A_i = 0. \end{aligned} \tag{5}$$

The RGE running of soft SUSY breaking masses as a function of the renormalization scale is shown in Fig. 1. After the RGE running, $m^2_{\tilde{N}^c_3}$ becomes negative while the other squared masses remain positive. The negative mass squared of the right-handed sneutrino triggers not only the $B-L$ symmetry breaking but also R-parity violation. Detailed analysis with random values of parameters has shown [10] that in most of the parameter region realizing the radiative $B-L$ symmetry breaking, R-parity is also broken.

We now analyze the scalar potential with the soft SUSY breaking parameters obtained from solving RGEs. Since the $B-L$ symmetry breaking scale is set to be 7 TeV in the following, we evaluate the RGE solutions at 7 TeV as follows:

$$\begin{aligned}
&g_{BL} = 0.286, \;\; y_1 = y_2 = 0.304, \; y_3 = 0.561, \\
&M_{BL} = 144 \text{ GeV}, \;\; m^2_{\tilde{N}^c_1} = m^2_{\tilde{N}^c_2} = 2.73 \times 10^6 \text{ GeV}^2, \;\; m^2_{\tilde{N}^c_3} = -3.57 \times 10^5 \text{ GeV}^2, \\
&m^2_{\Phi} = 6.03 \times 10^5 \text{ GeV}^2, \;\; m^2_{\bar{\Phi}} = 4.08 \times 10^6 \text{ GeV}^2, \\
&A_1 = A_2 = 34.1 \text{ GeV}, \;\; A_3 = 25.2 \text{ GeV}.
\end{aligned} \tag{6}$$

The scalar potential for $\tilde{N}^c_3$, Φ and $\bar{\Phi}$ consists of supersymmetric terms and soft SUSY breaking terms,

$$V = V_{SUSY} + V_{Soft}, \tag{7}$$

where

$$\begin{aligned}
V_{SUSY} &= |2y_3\tilde{N}^c_3\Phi|^2 + |\mu_\Phi\Phi|^2 + |y_3(\tilde{N}^c_3)^2 - \mu_\Phi\bar{\Phi}|^2 + \frac{g^2_{BL}}{2}\left(|\tilde{N}^c_3|^2 - 2|\Phi|^2 + 2|\bar{\Phi}|^2\right)^2, \\
V_{Soft} &= m^2_{\tilde{N}^c_3}|\tilde{N}^c_3|^2 + m^2_\Phi|\Phi|^2 + m^2_{\bar{\Phi}}|\bar{\Phi}|^2 - \left(A_3\Phi\tilde{N}^c_3\tilde{N}^c_3 + B_\Phi\bar{\Phi}\Phi + h.c.\right).
\end{aligned} \tag{8}$$

With appropriate values of μ_Φ and B_Φ, it is easy to numerically solve the stationary conditions for the scalar potential. For example, we find (in units of GeV)

$$\langle\tilde{N}^c_3\rangle = \frac{2762}{\sqrt{2}}, \;\; \langle\Phi\rangle = \frac{2108}{\sqrt{2}}, \;\; \langle\bar{\Phi}\rangle = \frac{2429}{\sqrt{2}} \tag{9}$$

for $\mu_\Phi = 2695$ GeV, $B_\Phi = 1.019 \times 10^7$ GeV^2 and the parameters given in Eq. (6). In this case, we have the Z' boson mass

$$m_{Z'} = g_{BL}v_{BL} = 2 \text{ TeV}, \tag{10}$$

where

$$v_{BL} = \sqrt{2\langle\tilde{N}^c_3\rangle^2 + 8\langle\Phi\rangle^2 + 8\langle\bar{\Phi}\rangle^2} = 7 \text{ TeV} \tag{11}$$

and the experimental lower bound $v_{BL} \geq 6 - 7$ TeV [14] is satisfied.

In order to prove that the stationary point is actually the potential minimum, we calculate the mass spectrum of the scalars, $\tilde{N}^c_3$, Φ and $\bar{\Phi}$. By straightforward numerical calculations, we find the eigenvalues of the mass matrix of the scalars $\Re[\tilde{N}^c_3]$, $\Re[\Phi]$ and $\Re[\bar{\Phi}]$ as $(1035, 1868, 5315)$ in GeV, while the mass eigenvalues for the pseudo-scalars $\Im[\tilde{N}^c_3]$, $\Im[\Phi]$ and $\Im[\bar{\Phi}]$ as $(0, 3308, 4858)$ in GeV. As expected, there is one massless eigenstate corresponding to the would-be Nambu-Goldstone mode. The other right-handed sneutrino mass eigenvalues are given by

$$\begin{aligned}
m^2_{\tilde{N}_{Ri}} &= m^2_{\tilde{N}^c_i} + 4y_i^2\langle\Phi\rangle^2 - 2y_iy_3\langle\tilde{N}^c_3\rangle^2 + 2A_i\langle\Phi\rangle + 2y_i\mu_\Phi\langle\bar{\Phi}\rangle + D_{BL}, \\
m^2_{\tilde{N}_{Ii}} &= m^2_{\tilde{N}^c_i} + 4y_i^2\langle\Phi\rangle^2 + 2y_iy_3\langle\tilde{N}^c_3\rangle^2 - 2A_i\langle\Phi\rangle - 2y_i\mu_\Phi\langle\bar{\Phi}\rangle + D_{BL},
\end{aligned} \tag{12}$$

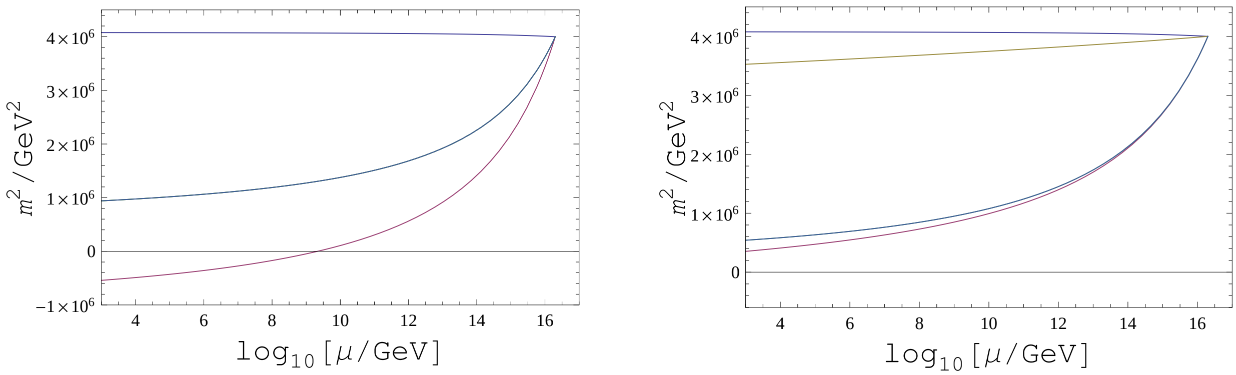

Figure 2: The RGE runnings of the soft SUSY breaking masses. The left panel shows the results for the boundary condition $y_1 = y_2 = y_3 = 1$ and others are the same as in Eq. (5), $m^2_{\bar{\Phi}}$, $m^2_{\tilde{N}_1^c} = m^2_{\tilde{N}_2^c} = m^2_{\tilde{N}_3^c}$ and m^2_{Φ} from above. The results for the boundary condition $y_1 = 0.2$, $y_2 = y_3 = 1$ are shown in the right panel, where $m^2_{\bar{\Phi}}$, $m^2_{\tilde{N}_1^c}$, $m^2_{\tilde{N}_2^c} = m^2_{\tilde{N}_3^c}$ and m^2_{Φ} from above.

where $m_{\tilde{N}_{Ri}}$ and $m_{\tilde{N}_{Ii}}$ $(i = 1, 2)$ are the mass eigenvalues for scalars and pseudo-scalars, respectively, and $D_{BL} = g^2_{BL}(\langle \tilde{N}_3^c \rangle^2 - 2\langle \Phi \rangle^2 + 2\langle \bar{\Phi} \rangle^2)$. We find $m_{\tilde{N}_{R1}} = m_{\tilde{N}_{R2}} = 2626$ GeV and $m_{\tilde{N}_{I1}} = m_{\tilde{N}_{I2}} = 1035$ GeV. Since the fermion components in $N^c_{2,3}$, Φ and $\bar{\Phi}$ and the $B-L$ gauginos are all mixed, it is quite involved to find the Majorana fermion mass eigenvalues. Accordingly, the seesaw mechanism is realized in a very complicated way. Although we do not discuss the fermion spectrum in detail here, our system with two right-handed neutrinos coupling to the SM neutrinos provides many free parameters; enough to reproduce the observed neutrino oscillation data. On the other hand, the mass of the Z_2-odd right-handed neutrino N_1^c is simply given by

$$M_{N_1^c} = 2y_1\langle \Phi \rangle = 906 \text{ GeV}. \quad (13)$$

Before concluding this section, it is interesting to observe the RGE running of the soft masses for different choices of the Majorana Yukawa couplings. The results are depicted in Fig. 2. In the left panel, we have taken the boundary condition $y_1 = y_2 = y_3 = 1$ but other parameters are the same as in Eq. (5), while the right panel shows the results for $y_1 = 0.2$, $y_2 = y_3 = 1$. In general, we can categorize the results of RGE runnings into three cases: (i) for only one $y_i = \mathcal{O}(1)$ and the others are relatively small, $m^2_{\tilde{N}_i^c}$ is driven to be negative, as a result, the $B-L$ symmetry as well as R-parity is broken at low energies. This case corresponds to Fig. 1. (ii) for $y_1 \sim y_2 \sim y_3 = \mathcal{O}(1)$, only m^2_{Φ} is driven to be negative. In this case, the $B-L$ symmetry is broken while R-party remains unbroken. This case corresponds to the left panel in Fig. 2. (iii) for $y_i \sim y_j = \mathcal{O}(1)$ $(i \neq j)$ and the other y_k $(k \neq i, j)$ is relatively small, all scalar squared masses remain positive and hence no symmetry breaking occurs. The right panel in Fig. 2 corresponds to this case.

4 Right-handed neutrino dark matter

As we showed in the previous section, the $B-L$ gauge symmetry is radiatively broken at the TeV scale. Associated with this radiative breaking, the right-handed sneutrino $\tilde{N}_3^c$ develops VEV and as a result, R-parity is also broken. Therefore, the neutralino is no longer the dark matter candidate. However, note that in our model the Z_2-parity is still exact, by which the lightest Z_2-odd particle is stable and can play the role of dark matter even in the presence of R-parity violation. As is evident in the mass spectrum we found in the previous section, the right-handed neutrino N_1^c is the lightest Z_2-odd particle. In this section, we evaluate the relic abundance of this right-handed neutrino dark matter candidate and identify the parameter region(s) consistent with the observations.

In [7], the relic abundance of the right-handed neutrino dark matter is analyzed in detail, where annihilation processes through the SM Higgs boson in the s-channel play the crucial role to reproduce the observed dark matter relic abundance. In the non-SUSY minimal $B-L$ model, the right-handed neutrino has a sizable coupling with the SM Higgs boson due to the mixing between the SM Higgs doublet and the $B-L$ Higgs in the scalar potential. However, in supersymmetric extension of the model there is no mixing between the MSSM Higgs doublets and the $B-L$ Higgs superfields in the starting superpotential. Although such a mixing emerges through the neutrino Dirac Yukawa coupling with the VEV of the right-handed sneutrino $\tilde{N}_3^c$, it is very small because of the small neutrino Dirac Yukawa coupling $y_D = \mathcal{O}(10^{-6})$. Among several annihilation channels of a pair of the Z_2-odd right-handed neutrinos, we find that the s-channel Z' boson exchange process gives the dominant contribution.

Now we evaluate the relic abundance of the right-handed neutrino by integrating the Boltzmann equation [19],

$$\frac{dY_{N_1^c}}{dx} = -\frac{x\gamma_{Z'}}{sH(M)}\left[\left(\frac{Y_{N_1^c}}{Y_{N_1^c}^{eq}}\right)^2 - 1\right], \tag{14}$$

where $Y_{N_1^c}$ is the yield (the ratio of the number density to the entropy density s) of the Z_2-odd right-handed neutrino, $Y_{N_1^c}^{eq}$ is the yield in thermal equilibrium, temperature of the universe is normalized by the mass of the right-handed neutrino $x = M/T$, and $H(M)$ is the Hubble parameter at $T = M$. The space-time densities of the scatterings mediated by the s-channel Z' boson exchange in thermal equilibrium are given by

$$\gamma_{Z'} = \frac{T}{64\pi^4}\int_{4M^2}^{\infty} ds\hat{\sigma}(s)\sqrt{s}K_1\left(\frac{\sqrt{s}}{T}\right), \tag{15}$$

where s is the squared center-of-mass energy, K_1 are the modified Bessel function of the first kind, and the total reduced cross section for the process $N_1^c N_1^c \to Z' \to f\bar{f}$ (f denotes the SM fermions) is

$$\hat{\sigma}_{Z'}(s) = \frac{26}{12\pi} g_{BL}^4 \frac{\sqrt{s}\,(s - 4M^2)^{\frac{3}{2}}}{\left(s - m_{Z'}^2\right)^2 + m_{Z'}^2 \Gamma_{Z'}^2} \tag{16}$$

with the decay width of the Z' boson,

$$\Gamma_{Z'} = \frac{g_{BL}^2}{24\pi} \left[13 + 2 \left(1 - \frac{4M^2}{m_{Z'}^2} \right)^{\frac{3}{2}} \theta \left(m_{Z'}^2 / M^2 - 4 \right) \right]. \tag{17}$$

For simplicity, we have assumed that $y_1 = y_2$ as in the previous section and that the other particles (except for the SM particles) are all heavy with mass $> m_{Z'}/2$. This assumption is consistent with the parameter choice in our analysis below.

Now we solve the Boltzmann equation numerically. To solve the equation for the relevant domain, we inherit parameter values from those presented in the previous section which were already motivated as interesting values,

$$g_{BL} = 0.286, \quad m_{Z'} = 2 \text{ TeV}, \tag{18}$$

while $M_{N_1^c} = M$ is taken to be a free parameter. With the asymptotic value of the yield $Y_{N_1^c}(\infty)$ the dark matter relic density is written as

$$\Omega h^2 = \frac{M s_0 Y_{N_1^c}(\infty)}{\rho_c / h^2}, \tag{19}$$

where $s_0 = 2890$ cm^{-3} is the entropy density of the present universe, and $\rho_c/h^2 = 1.05 \times 10^{-5}$ GeV/cm^3 is the critical density. The result should be compared with the observations at 2σ level [20]

$$\Omega_{DM} h^2 = 0.1120 \pm 0.0056. \tag{20}$$

Fig. 3 shows the relic abundance of the right-handed neutrino dark matter as a function of its mass. The dashed lines correspond to the upper and the lower bounds on the dark matter relic abundance in Eq. (20). We find two solutions

$$M \simeq 906, \quad 1016 \text{ GeV}. \tag{21}$$

It turns out from Fig. 3 that in order to reproduce the observed relic abundance, the enhancement of the annihilation cross section is necessary, so that the mass of the dark matter should be close to the Z' boson resonance point. The dark matter mass $M = 906$ GeV coincides with the value presented in the previous section. For a different parameter choice, the Z_2-odd right-handed sneutrino (scalar or pseudo-scalar) can be the lightest Z_2-odd particle and a candidate for the dark matter, instead of the right-handed neutrino. However, in this case, there is no s-channel Z' boson mediated process and the resultant relic abundance is found to be too large.

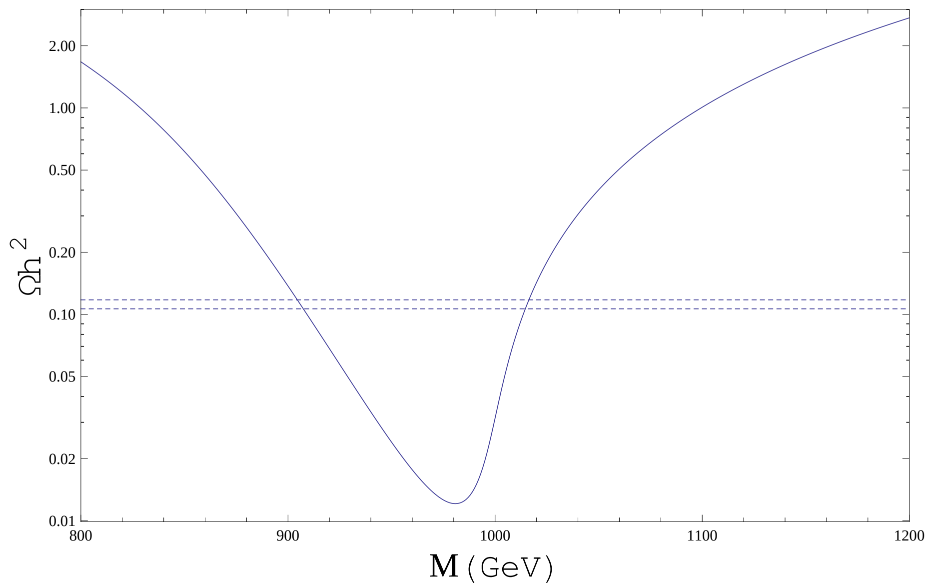

Figure 3: The relic abundance of the dark matter right-handed neutrino as a function of its mass. The dashed lines represent the upper and the lower bounds on the dark matter relic abundance.

5 Phenomenological constraints

The $B-L$ symmetry works to forbid phenomenologically dangerous R-parity violating terms in the MSSM. Although this is a remarkable advantage of the $B-L$ extended model, the $B-L$ symmetry is eventually broken at the TeV scale, in particular, R-parity violation also occurs by non-zero VEV of the right-handed sneutrino $\tilde{N}_3^c$. We here consider phenomenological constraints on our model associated with the R-parity violation.

Once $\langle \tilde{N}_3^c \rangle \neq 0$ has developed the so-called bilinear R-party violating term is generated in Eq. (1),

$$W_{BL} \supset \sum_{j=1}^{3} y_D^{3j} \langle \tilde{N}_3^c \rangle L_j H_u. \tag{22}$$

These terms induce the lepton number violating Yukawa couplings in the MSSM superpotential,

$$W_{\Delta L=1} \sim \left(y_e^{ij} \epsilon^k \right) E_i^c L_j L_k + \left(y_d^{ij} \epsilon^k \right) D_i^c Q_j L_k, \tag{23}$$

where y_e and y_d are Yukawa matrices for the charged leptons and the down-type quarks, and

$$\epsilon^k = y_D^{3k} \frac{\langle \tilde{N}_3^c \rangle}{\mu} \tag{24}$$

is the mixing parameters between H_d and L_i with the μ-parameter for the Higgs doublets in the MSSM. Through the interactions, lepton-number-changing processes can be active in the early universe and erase an existing baryon asymmetry. The requirement that this erasure should not occur before the electroweak transition typically gives [21]

$$\left(y_e^{ij}\epsilon^k\right),\ \left(y_d^{ij}\epsilon^k\right) \lesssim 10^{-7}, \tag{25}$$

which in turn implies

$$\epsilon^k \lesssim 10^{-5}\cos\beta. \tag{26}$$

This condition is marginally satisfied with the natural order of magnitude for $y_D \sim 10^{-6}$ in the seesaw mechanism at the TeV scale. Note that this is a sufficient condition, and some flavor structures can relax it [21].

It would be fair to mention that although the renormalizable R-parity violating terms in the MSSM are forbidden by the $B-L$ symmetry, one may introduce dimension five $B-L$ symmetric operators such as $\kappa_{ijk\ell}Q_iQ_jQ_kL_\ell/M_P$. Even though these are Planck-scale suppressed operators, they can lead to rapid proton decay. In fact, experimental limits from nucleon stability are very severe, $\kappa_{1121}, \kappa_{1122} \lesssim 10^{-8}$ [22]. We need to assume such small parameters. Since the above dimension five operator is also R-parity symmetric, the same discussion is applied to the MSSM.

6 Conclusions and discussions

The minimal gauged $\mathrm{U}(1)_{B-L}$ model based on the gauge group $\mathrm{SU}(3)_c\times\mathrm{SU}(2)_L\times\mathrm{U}(1)_Y\times\mathrm{U}(1)_{B-L}$ is an elegant and simple extension of the Standard Model, in which the right-handed neutrinos of three generations are necessarily introduced for the gauge and gravitational anomaly cancellations. The mass of right-handed neutrinos arises associated with the $\mathrm{U}(1)_{B-L}$ gauge symmetry breaking, and the seesaw mechanism is naturally implemented. Supersymmetric extension of the minimal $B-L$ model offers not only a solution to the gauge hierarchy problem but also a natural mechanism of breaking the $B-L$ symmetry at the TeV scale, namely, the radiative $B-L$ symmetry breaking. Although the radiative symmetry breaking at the TeV scale is a remarkable feature of the model, R-parity is also broken by non-zero VEV of a right-handed sneutrino. Therefore, the neutralino, which is the conventional dark matter candidate in SUSY models, becomes unstable and cannot play the role of the dark matter any more.

We have proposed introducing a Z_2-parity and assigned an odd-parity to one right-handed neutrino. This parity ensures the stability of the right-handed neutrino and hence the right-handed neutrino can be a dark matter candidate even in the presence of R-parity violation. No new particle introduced for the dark matter is required. We have shown that for a parameter set, the mass squared of a right-handed sneutrino is driven to be negative by the RGE running. Analyzing the scalar potential with RGE solutions of soft SUSY breaking parameters, we have identified the vacuum where the $B-L$ symmetry as well as R-parity is broken at the TeV scale.

We have numerically integrate the Boltzmann equation for the Z_2-odd right-handed neutrino and found that its relic abundance is consistent with the observations. In reproducing the observed dark matter relic density, an enhancement of the annihilation cross section via the Z' boson s-channel resonance is necessary, so that the dark matter mass should be close to half of Z' boson mass.

Associated with the $B-L$ symmetry breaking, all new particles have TeV-scale masses, which is being tested at the LHC in operation. Discovery of the Z' boson resonance at the LHC [23] is the first step to confirm our model. Once the Z' boson mass is measured, the dark matter mass is also determined in our model. If kinematically allowed, the Z' boson decays to the dark matter particles with the branching ratio ~ 0.6 % (see Eq. (17)). Precise measurement of the invisible decay width of Z' boson can reveal the existence of the dark matter particle.

A variety of experiments are underway to directly or indirectly detect dark matter particles. For the detection, it is crucial for a dark matter particle to have sizable spin-independent and/or spin-dependent elastic scattering cross sections with nuclei. In our model, the right-handed neutrino dark matter couples with quarks in two ways. One is via the Z' boson exchange process, the other is MSSM Higgs boson mediated processes. Because of its Majorana nature, the dark matter particle has the axial vector coupling with the Z' boson, while the SM fermions have the vector couplings. As a result, the Z' boson exchange process has no contribution to the elastic scattering between the dark matter particle and quarks in the non-relativistic limit. Although the right-handed neutrino dark matter has no direct coupling with the MSSM Higgs bosons, such couplings are generated after the $B-L$ and electroweak symmetry breakings. For example, a mixing mass, $|y_3|^2\langle\Phi\rangle\langle\tilde{N}_3^c\rangle$, between Φ and $\tilde{N}_3^c$ is generated by the radiative $B-L$ symmetry breaking and after the electroweak symmetry breaking, another mixing mass such as $|y_D|^2\langle\tilde{N}_3^c\rangle\langle H_u\rangle$ between $\tilde{N}_3^c$ and H_u is generated. These mixing masses eventually induce the couplings between the right-handed neutrino dark matter and the MSSM Higgs bosons. However, the coupling constants are suppressed by the small factor $|y_D|^2 \sim 10^{-12}$, so that the spin-independent cross section for the elastic scattering off nuclei mediated by the MSSM Higgs boson exchanges is too small to be detected.

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