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Vectorlike sneutrino dark matter

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In this paper, we discuss the minimal supersymmetric standard model (MSSM) extended with one vectorlike lepton doublet $L-\bar{L}$ and one right-handed neutrino N. The neutral vecotorlike sneutrino can be a candidate of dark matter. To avoid the interaction with the nucleons by exchanging a Z boson, the mass splitting between the real part and the imaginary part of the sneutrino field is needed. Compared with the MSSM sneutrino dark matter, the mass splitting between the vectorlike sneutrino field can be more naturally acquired without large A terms and constraints on the neutralino masses. We have also calculated the relic density and the elastic scattering cross sections with the nucleons in the cases that the dark matter particles coannihilate with or without the MSSM slepton doublets. The elastic scattering cross sections with the nucleons are well below the LUX bounds. In the case that the dark matter coannihilates with all the MSSM slepton doublets, the mass of the dark matter can be as light as 370 GeV.

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I. INTRODUCTION

In the supersymmetric models, R parity $(-1)^{(3B+L+2S)}$ usually conserves in order to forbid the protons to decay (for a review, see Ref. [1]). Then, the lightest supersymmetric particle (LSP) can become dark matter if it is neutral. Neutralinos and sneutrinos have been considered as the candidates of dark matter in the literature. However, compared with the neutralinos, sneutrinos in the minimal supersymmetric standard model (MSSM) suffer from the difficulty in escaping the direct detection bounds since they can exchange a Z boson with the nucleons [2]. One way to avoid this problem is to introduce the mass splitting between the real part and the imaginary part of the sneutrino field [3–8]. This trick has been applied in many inelastic dark matter models (for examples, see Refs. [9–13]). To achieve this splitting, we need some lepton number-violating sectors beyond the MSSM, which would arise from either the right-handed neutrinos or some $SU(2)_L$-triplet Higgs fields. These sectors can also make up for the deficiency of the MSSM that the neutrinos are massless. However, in order to acquire enough splitting value $|m_{\tilde{\nu}^+}-m_{\tilde{\nu}^-}|\gtrsim 100$ KeV and at the same time keep the sub-eV masses of the light neutrinos, large A terms are usually required, and limits on the masses of the neutralinos are also imposed.

In this paper, we discuss a model that extends the MSSM with a pair of vectorlike leptons $(L+\bar{L})$. If the vectorlike sneutrinos end up as the dark matter, we also need to split the real part and the imaginary part of the vectorlike sneutrino field. The simplest way to achieve this is to introduce another right-handed neutrino field N together with the lepton number-violating terms motivated from the type I seesaw mechanisms [14–18]. We will see that in this model enough mass splitting can arise from the LH_uN and $\bar{L}H_dN$ Yukawa terms even if we switch off all the trilinear A terms. The values of these Yukawa coupling constants can have an impact on the relic density of the dark matter and can also contribute to the direct detection signals. If the mixings between the vectorlike sectors and the MSSM sectors are small enough, the sub-eV neutrino masses can also remain undisturbed, relaxing the bounds on the masses of the neutralinos. In the literature, there are models in which the MSSM is extended with the vectorlike particles (for examples, see Refs. [19–28]). Vectorlike sectors can either be heavier than the 100 TeV scale and play the role of so-called "messengers" in the gauge-mediating supersymmetry breaking (GMSB) models or influence the TeV-scale phenomenologies if the vectorlike particles are relatively light. The latter case is particularly interesting partly because TeV-scale vectorlike particles can be tested directly through collider searches in the LHC era. Vectorlike particles can also interact with the Higgs sectors, relieving the little-hierarchy problem to reach the sufficient standard model (SM)-like Higgs mass in the MSSM.

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We should note that in order to keep the unification of the gauge-coupling constants our model can be embedded in a $5+\bar{5}$ model, which also contains a pair of vectorlike down-type quarks ($D+\bar{D}$). However, in the following text, we disregard this. In Refs. [29,30], there is a similar model in which the vectorlike messenger sleptons as light as 1 to 3 TeV play the role of the dark matter in the framework of the GMSB models (for a review, see Ref. [31]). However, in this paper, we do not concern the origin of the breaking of the supersymmetry, and the vectorlike leptons just sense the supersymmetry breaking indirectly, just similar to the ordinary MSSM fields. The dark matter can become much lighter when coannihilating with the MSSM sleptons in our model.

This paper is organized as follows. Section II describes the model, and calculations of the mass matrices are presented. Section III calculates the relic density and the spin-independent cross section with the nucleons numerically. The Yukawa couplings constants are adjusted for a best fit to the Planck result of relic density [32]. Finally, Sec. IV contains the conclusions and discussions.

II. MODEL DESCRIPTIONS

Besides the MSSM chiral superfields H_u, H_d, L_i, E_i, Q_i, U_i, and D_i ($i=1$–3), which are the up-type Higgs doublet, down-type Higgs doublet, the left-handed lepton doublets, the right-handed charged leptons, the left-handed quark doublets, and the up-type and the down-type right-handed quarks of the three generations, respectively, we introduce L, $\bar{L}$, N in our model, which are a pair of vectorlike lepton doublets and one right-handed neutrino. They are assigned with the odd R parity. The involved superpotential is given by

$$W \supset \mu_L L\bar{L} + y_L L H_u N + y_{\bar{L}} \bar{L} H_d N + \mu_N N^2 + \mu H_u H_d. \quad (1)$$

The supersymmetric breaking soft mass terms and the trilinear A terms are given by

$$\begin{aligned}\mathcal{L}_{\text{soft}} \supset\, & m_L^2|\tilde{L}|^2 + m_{\bar{L}}^2|\tilde{\bar{L}}|^2 + m_N^2|\tilde{N}|^2 + B_N\mu_N(\tilde{N}^2 + \text{H.c.}) \\ & + B_L\mu_L(\tilde{L}\,\tilde{\bar{L}} + \text{H.c.}) + (A_{y_L} y_L \tilde{L} H_u \tilde{N} \\ & + A_{y_{\bar{L}}} y_{\bar{L}} \tilde{\bar{L}} H_d \tilde{N} + \text{H.c.}). \end{aligned} \quad (2)$$

Generally speaking, Eqs. (1) and (2) do not contain all possible terms which conserve the $U(1)_Y \times SU(2)_L \times SU(3)_C$ quantum numbers and the R parity. These terms can result in the mixings between the MSSM sectors and the vectorlike sectors (e.g., $\tilde{L}_i^\dagger \tilde{L}$), and can also lead to the light-neutrino mass through both the tree-level Type I seesaw mechanisms and loop-level effects [33,34] (e.g., $y_i L_i H_u N$, together with the corresponding A terms). In the former case, we assume these terms are small enough to be omitted not only for simplicity but also because of the precision electroweak constraints on the mixings between the MSSM and the vectorlike sectors. For the latter case, the detailed specific mass spectrum and the mixing patterns of the neutrino sectors are out of the scope of this paper, and the smallness of the neutrino masses suppresses the effects from these terms. However, we should note that all these terms cannot be totally absent, because in some coannihilation cases to be discussed, these terms supply the way for the coannihilating particles to finally decay into dark matter particles.

The conventions of the vacuum expectation values (VEVs) of the Higgs sectors are

$$H_u^0 = v_u + \frac{R_u + iI_u}{\sqrt{2}}, \qquad H_d^0 = v_d + \frac{R_d + iI_d}{\sqrt{2}}. \quad (3)$$

After the Higgs doublets acquire the VEVs, the real part and the imaginary part of the vectorlike neutral sneutrinos are separated. We define

$$\tilde{L} = \begin{bmatrix} \frac{R_L + iI_L}{\sqrt{2}} \\ \tilde{L}^- \end{bmatrix}, \quad \tilde{\bar{L}} = \begin{bmatrix} \tilde{\bar{L}}^+ \\ \frac{R_{\bar{L}} + iI_{\bar{L}}}{\sqrt{2}} \end{bmatrix}, \quad N = \frac{R_N + iI_N}{\sqrt{2}}. \quad (4)$$

The mass matrices are therefore

$$V \supset \frac{1}{2}[R_L, R_{\bar{L}}, R_N]\mathcal{M}_R \begin{bmatrix} R_L \\ R_{\bar{L}} \\ R_N \end{bmatrix} + \frac{1}{2}[I_L, I_{\bar{L}}, I_N]\mathcal{M}_I \begin{bmatrix} I_L \\ I_{\bar{L}} \\ I_N \end{bmatrix}, \quad (5)$$

where

$$\begin{aligned} \mathcal{M}_R &= \mathcal{M}_{RF} + \mathcal{M}_{RD} + \mathcal{M}_{RS}, \\ \mathcal{M}_I &= \mathcal{M}_{IF} + \mathcal{M}_{ID} + \mathcal{M}_{IS}. \end{aligned} \quad (6)$$

The matrix elements originating from the F terms are

$$\mathcal{M}_{RF} = \begin{bmatrix} y_L^2 v_u^2 + \mu_L^2 & -y_L y_{\bar{L}} v_u v_d & -y_L \mu v_d - y_{\bar{L}} v_d \mu_L + 2y_L v_u \mu_N \\ -y_L y_{\bar{L}} v_u v_d & y_{\bar{L}}^2 v_d^2 + \mu_L^2 & y_{\bar{L}} v_u \mu + y_L \mu_L v_u - 2y_{\bar{L}} v_d \mu_N \\ -y_L v_d \mu - y_{\bar{L}} v_d \mu_L + 2y_L \mu_N v_u & y_{\bar{L}} v_u \mu + y_L v_u \mu_L - 2y_{\bar{L}} v_d \mu_N & y_{\bar{L}}^2 v_u^2 + y_{\bar{L}}^2 v_d^2 + 4\mu_N^2 \end{bmatrix},$$

$$\mathcal{M}_{IF} = \begin{bmatrix} y_L^2 v_u^2 + \mu_L^2 & -y_L y_{\bar{L}} v_u v_d & y_L v_d \mu - y_{\bar{L}} v_d \mu_L + 2y_L v_u \mu_N \\ -y_L y_{\bar{L}} v_u v_d & y_{\bar{L}}^2 v_d^2 + \mu_L^2 & -y_{\bar{L}} v_u \mu + y_L v_u \mu_L - 2y_{\bar{L}} v_d \mu_N \\ y_L v_d \mu - y_{\bar{L}} v_d \mu_L + 2y_L v_u \mu_N & -y_{\bar{L}} v_u \mu + y_L v_u \mu_L - 2y_{\bar{L}} v_d \mu_N & y_{\bar{L}}^2 v_u^2 + y_{\bar{L}}^2 v_d^2 + 4\mu_N^2 \end{bmatrix}. \quad (7)$$

The matrix elements induced by the gauge D terms are

$$\mathcal{M}_{RD,11} = \mathcal{M}_{ID,11} = \frac{1}{4}(-g_1^2 v_u^2 + g_1^2 v_d^2 - g_2^2 v_u^2 + g_2^2 v_d^2),$$
$$\mathcal{M}_{RD,22} = \mathcal{M}_{ID,22} = \frac{1}{4}(g_1^2 v_u^2 - g_1^2 v_d^2 + g_2^2 v_u^2 - g_2^2 v_d^2), \quad (8)$$

and all the other matrix elements of the $\mathcal{M}_{RD}$ and the $\mathcal{M}_{ID}$ equal zero. $g_{1,2}$ are the $U(1)_Y$ and the $SU(2)_L$ gauge coupling constants, respectively. The matrix elements induced by the soft terms are

$$\mathcal{M}_{RS} = \begin{bmatrix} m_L^2 & B_L\mu_L & y_L A_{y_L} v_u \\ B_L\mu_L & m_{\bar{L}}^2 & y_{\bar{L}} A_{y_{\bar{L}}} v_d \\ y_L A_{y_L} v_u & y_{\bar{L}} A_{y_{\bar{L}}} v_d & m_N^2 + B_N\mu_N \end{bmatrix},$$
$$\mathcal{M}_{IS} = \begin{bmatrix} m_L^2 & -B_L\mu_L & -y_L A_{y_L} v_u \\ -B_L\mu_L & m_{\bar{L}}^2 & -y_{\bar{L}} A_{y_{\bar{L}}} v_d \\ -y_L A_{y_L} v_u & -y_{\bar{L}} A_{y_{\bar{L}}} v_d & m_N^2 - B_N\mu_N \end{bmatrix}. \quad (9)$$

After diagonalizing $\mathcal{M}_{R,I}$, we acquire three CP-even and CP-odd real scalar particles $R_{1,2,3}$ and $I_{1,2,3}$. They are defined as

$$R_L = Z_{R11}R_1 + Z_{R12}R_2 + Z_{R13}R_3,$$
$$R_{\bar{L}} = Z_{R21}R_1 + Z_{R22}R_2 + Z_{R23}R_3,$$
$$R_N = Z_{R31}R_1 + Z_{R32}R_2 + Z_{R33}R_3,$$
$$I_L = Z_{I11}I_1 + Z_{I12}I_2 + Z_{I13}I_3,$$
$$I_{\bar{L}} = Z_{I21}I_1 + Z_{22}I_2 + Z_{23}I_3,$$
$$I_N = Z_{I31}I_1 + Z_{I32}I_2 + Z_{33}I_3, \quad (10)$$

where $Z_{I,Rij}$'s are the matrix elements of the diagonalizing matrices. Without loss of generality, we assign an ascending order of masses among $R_{1,2,3}$ and $I_{1,2,3}$. The mass matrix of the charged vectorlike sleptons is

$$V \supset [\tilde{L}^{-*}, \tilde{\bar{L}}^{+}]\mathcal{M}_{\tilde{L}^\pm}\begin{bmatrix} \tilde{L}^- \\ \tilde{\bar{L}}^{+*} \end{bmatrix}, \quad (11)$$

where

$$\mathcal{M}_{\tilde{L}^\pm} = \mathcal{M}_{\tilde{L}^\pm F} + \mathcal{M}_{\tilde{L}^\pm D} + \mathcal{M}_{\tilde{L}^\pm S}. \quad (12)$$

The elements originating from the F terms are simply

$$\mathcal{M}_{\tilde{L}^\pm F11} = \mathcal{M}_{\tilde{L}^\pm F11} = \mu_L^2, \quad \mathcal{M}_{\tilde{L}^\pm F12} = \mathcal{M}_{\tilde{L}^\pm F21} = 0. \quad (13)$$

The elements induced by the D terms are

$$\mathcal{M}_{\tilde{L}^\pm D11} = \frac{1}{4}g_1^2 v_d^2 - \frac{1}{4}g_2^2 v_d^2 - \frac{1}{4}g_1^2 v_u^2 + \frac{1}{4}g_2^2 v_u^2$$
$$\mathcal{M}_{\tilde{L}^\pm D22} = -\frac{1}{4}g_1^2 v_d^2 + \frac{1}{4}g_2^2 v_d^2 + \frac{1}{4}g_1^2 v_u^2 - \frac{1}{4}g_2^2 v_u^2$$
$$\mathcal{M}_{\tilde{L}^\pm D12} = \mathcal{M}_{\tilde{L}^\pm D21} = 0. \quad (14)$$

The matrix elements induced by the soft terms are

$$\mathcal{M}_{\tilde{L}^\pm S11} = \begin{bmatrix} m_L^2 & -B_L\mu_L \\ -B_L\mu_L & m_{\bar{L}}^2 \end{bmatrix}. \quad (15)$$

After diagonalizing the $\mathcal{M}_{L^\pm}$, we acquire two charged sleptons,

$$\tilde{L}^- = Z_{c11}\tilde{L}_1^- + Z_{c12}\tilde{L}_2^-, \quad \tilde{\bar{L}}^{+*} = Z_{c21}\tilde{L}_1^- + Z_{c22}\tilde{L}_2^-, \quad (16)$$

where Z_{cij}'s are the diagonalizing matrix elements. The mass matrix of the vectorlike neutrinos together with the right-handed neutrino is given by

$$\mathcal{L} \supset \frac{1}{2}[L^{0C}, \bar{L}^{0C}, N^C]\mathcal{M}_{L^0}\begin{bmatrix} L^0 \\ \bar{L}^0 \\ N \end{bmatrix}, \quad (17)$$

where $X^c = X^\dagger \cdot (i\sigma^2)$, σ^i $(i = 1,\ 2,\ 3)$ are the Pauli matrices, and X is a two-component Weyl spinor. The matrix elements of the $\mathcal{M}_{L^0}$ are

$$\mathcal{M}_{L^0} = \begin{bmatrix} 0 & \mu_L & y_L v_u \\ \mu_L & 0 & -y_{\bar{L}} v_d \\ y_L v_u & -y_{\bar{L}} v_d & 2\mu_N \end{bmatrix}. \quad (18)$$

After diagonalizing the $\mathcal{M}_{L^0}$, we acquire these three neutral majorana fermions,

$$L^0 = Z_{011}L_1^0 + Z_{012}L_2^0 + Z_{013}L_3^0,$$
$$\bar{L}^0 = Z_{021}L_1^0 + Z_{022}L_2^0 + Z_{023}L_3^0,$$
$$N = Z_{031}L_1^0 + Z_{032}L_2^0 + Z_{033}L_3^0, \quad (19)$$

where Z_{0ij}'s are the diagonalizing matrix elements.

Finally, L^- and $\bar{L}^+$ form a Dirac fermion, and its mass is μ_L.

From observing (7), we can learn that, although $\mathcal{M}_{RF,11} = \mathcal{M}_{IF,11}$, $\mathcal{M}_{RF,22} = \mathcal{M}_{IF,22}$, the off-diagonal $|\mathcal{M}_{RF,13}| \neq |\mathcal{M}_{IF,13}|$. This will split the mass between the R_i's and I_i's even if we switch off all the mass terms induced by the D terms and the A terms. In some cases, this difference can be well estimated. For example, if $m_{\bar{L}}^2$, $\mu_N^2 \gg m_L^2$, the lightest two scalar fields, say R_1 and I_1, would be dominated by R_L and I_L; then,

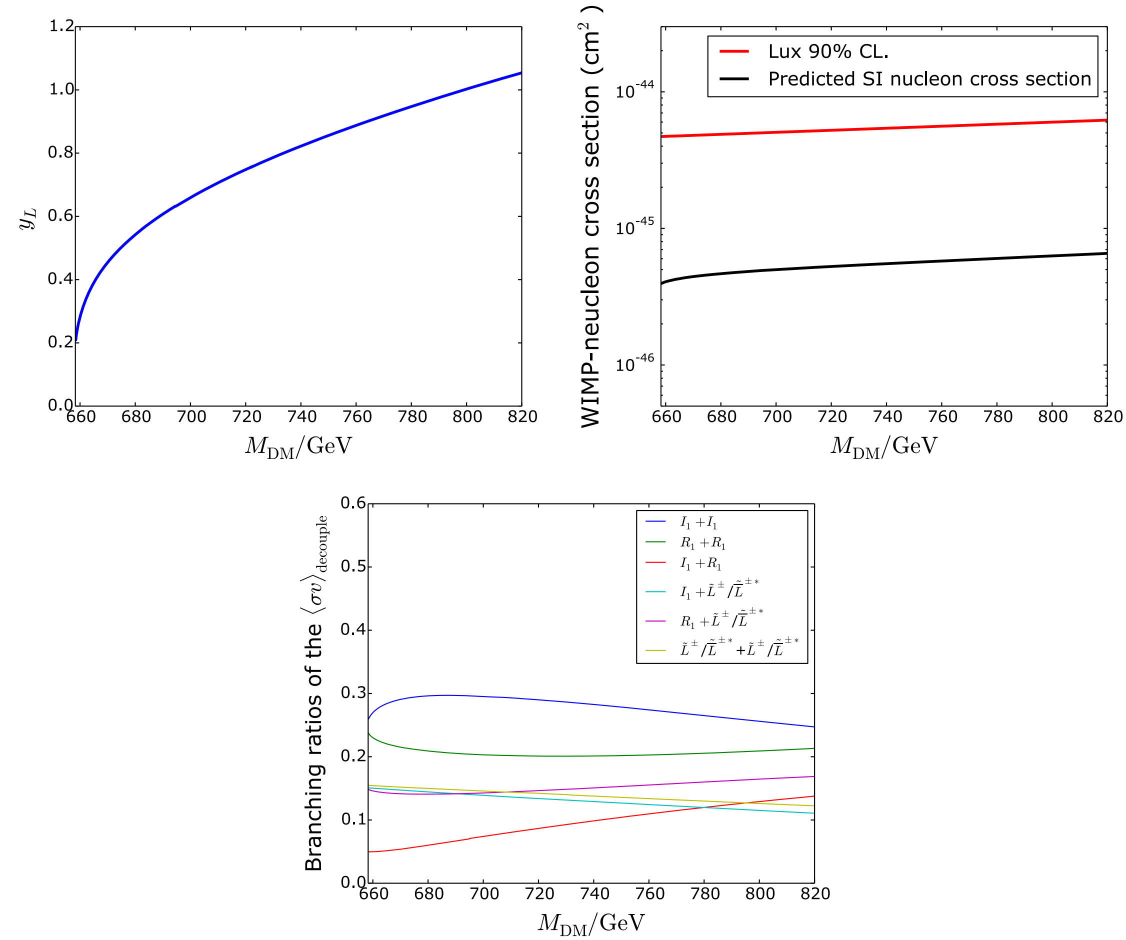

FIG. 1. The y_L corresponding to $\Omega_c h^2 = 0.1199$ (left panel), the spin-independent cross section with the nucleons of the dark matter particles (right panel), and the branching ratios of $\langle\sigma v\rangle_{\mathrm{decouple}}$ (bottom panel) in the case in which only the I_1, R_1, together with $\tilde{L}_1$ coannihilate.

$$m_{R_1}^2 - m_{I_1}^2 \approx -\frac{(-y_L\mu v_d - y_{\bar{L}} v_d\mu_L + 2y_L v_u\mu_N)^2}{4\mu_N^2} + \frac{(y_L\mu v_d - y_{\bar{L}} v_d\mu_L + 2y_L v_u\mu_N)^2}{4\mu_N^2} = \frac{2y_L^2\mu v_d v_u\mu_N - y_L y_{\bar{L}} v_d^2\mu\mu_L}{\mu_N^2}, \quad (20)$$

so

$$m_{R_1} - m_{I_1} \approx \frac{1}{\overline{m_{R,I_1}}}\frac{y_L^2\mu v_d v_u}{\mu_N}, \quad (21)$$

where $\overline{m_{R,I_1}}$ is the average value of the masses of R_1 and I_1. For example, if $\mu = 500$ GeV, $\tan\beta = \frac{v_u}{v_d} = 15$, $\mu_N = 1$ TeV, $y_{\bar{L}} = y_L = 0.1$, and $\overline{m_{R,I_1}} = 400$ GeV, then $m_{R_1} - m_{I_1} \approx 20$ MeV, which is far beyond the needed $O(100\ \mathrm{KeV})$ in order to escape the direct detection bounds. In this scenario, I_1 will be lighter than R_1, which means I_1 tend to become the dark matter if all the $A_{y_L,y_{\bar{L}}}$, $B_{L,N}$ terms are set zero.

III. NUMERICAL RESULTS OF RELIC ABUNDANCE AND DIRECT DETECTION

If $m_L^2 \approx m_{\bar{L}}^2$, the masses of the $I_{L,\bar{L}}$, $R_{L,\bar{L}}$, $\tilde{L}^-$, $\tilde{\bar{L}}^+$ are close to each other, and there are large mixings between the neutral and the charged sleptons, respectively. For a clearer aspect, we assume a large difference between the m_L^2 and the $m_{\bar{L}}^2$ in this paper to avoid the rather complicated mixings and coannihilating cases. The right-handed (s)neutrino mass terms m_N^2, μ_N are also large enough for the right-handed (s)neutrinos to decouple during the annihilating processes. In this situation, the mixings between the right-handed sneutrinos and the vectorlike sneutrinos are also suppressed by their large mass differences.

According to Eqs. (8) and (14), the mass terms induced by the D-terms lower the masses of the R, I_L-dominated particle and increase the mass of the L^--dominated charged sneutrino, while these terms lower the masses of the R, $I_{\bar{L}}$-dominated particle and give rise to the mass of the $\bar{L}^{+*}$-dominated charged sneutrino. It means that if $m_{\bar{L}}^2 \ll m_L^2$ the masses of the R, $I_{\bar{L}}$-dominated particles tend to be a little

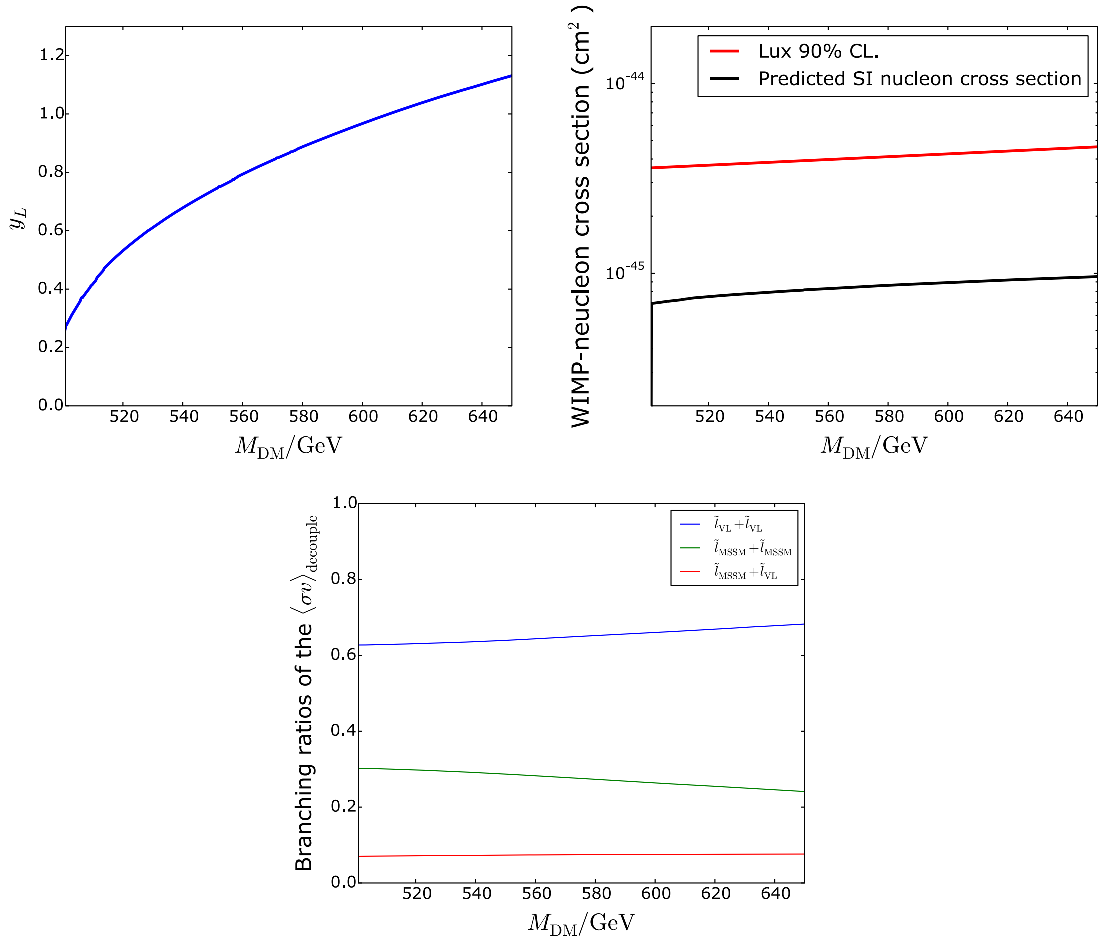

FIG. 2. The y_L corresponding to $\Omega_c h^2 = 0.1199$ (left panel), the spin-independent cross section with the nucleons of the dark matter particles (right panel), and the branching ratios of $\langle\sigma v\rangle_{\rm decouple}$ (bottom panel) in the case in which the vectorlike sleptons coannihilate with one generation of the MSSM slepton.

heavier than the charged $\bar{L}^{+*}$-dominated particle, leaving us a charged LSP in most cases. Because of this, we assume $m_{\tilde{L}}^2 \ll m_{\tilde{\bar{L}}}^2$ in the following text. As was discussed in the previous section, it means that the LSP will be a CP-odd I_L-dominated I_1.

A terms also play roles in the annihilating processes. However, as we have noted, $\tilde{N}$ decouples, so both the effects from the $A_{y_L} y_L \tilde{L} H_u \tilde{N}$ and the $A_{y_{\bar{L}}} y_{\bar{L}} \tilde{\bar{L}} H_d \tilde{N}$ terms are suppressed. Although A terms also modify the mass spectrum of the supersymmetric particles, numerical calculations also show that $A_{y_L, y_{\bar{L}}} \sim O(100\ \text{GeV})$ does not influence the final results to a notable extent. According to all these reasons, we set $A_{y_L} = A_{y_{\bar{L}}} = 0$ in the following discussions.

For simplicity, we also assume that all the other MSSM sparticles and the exotic Higgs bosons decouple except the Binos ($\tilde{B}$), Winos ($\tilde{W}^{\pm,0}$), and some $SU(2)_L$ doublet sleptons in some coannihilating cases. We set the masses of all the Binos and Winos to be $m_{\tilde{B}} = m_{\tilde{W}^{\pm,0}} = 2$ TeV. We also set the alignment condition $\beta = \frac{\pi}{2} - \alpha$, where α is the neutral Higgs bosons' mixing angle. This equals the $m_A \to \infty$ limit, where m_A is the mass of the CP-odd Higgs boson. We set $\mu_L = 300$ GeV during the calculation, which is safe from the bounds on heavy leptons [35].

The model is implemented with the FeynRules 2.3.12 [36] to generate the CalcHEP [37] model files. Then, MicrOMEGAs 4.2.5 [38] is used to calculate the relic density, the spin-independent cross section with the nucleons, and the branching ratios contributing to the $\langle\sigma v\rangle_{\rm decouple}$, which is the annihilation cross section between the dark matter particles when they decouple. For each mass of the dark matter, we calculate the y_L which corresponds to the best-fit Planck data $\Omega_c h^2 = 0.1199$ [32] and plot the $m_{\rm DM}$, y_L, branching ratios contributing to the $\langle\sigma v\rangle_{\rm decouple}$ and the spin-independent direct detection cross section with the nucleon $\sigma_{\rm SI}$ in four cases, which are no coannihilation, coannihilation with one MSSM slepton, coannihilations with two MSSM sleptons, and coannihilations with three MSSM sleptons in Figs. 1, 2, 3, and 4. For each coannihilating situation, we guarantee the masses of the coannihilating MSSM sneutrinos to be 2 GeV heavier than the mass of the dark matter. Note that it is impossible and unnecessary to plot every branching ratio of the $\langle\sigma v\rangle_{\rm decouple}$ in such small graphs, so we sum over the channels according to the

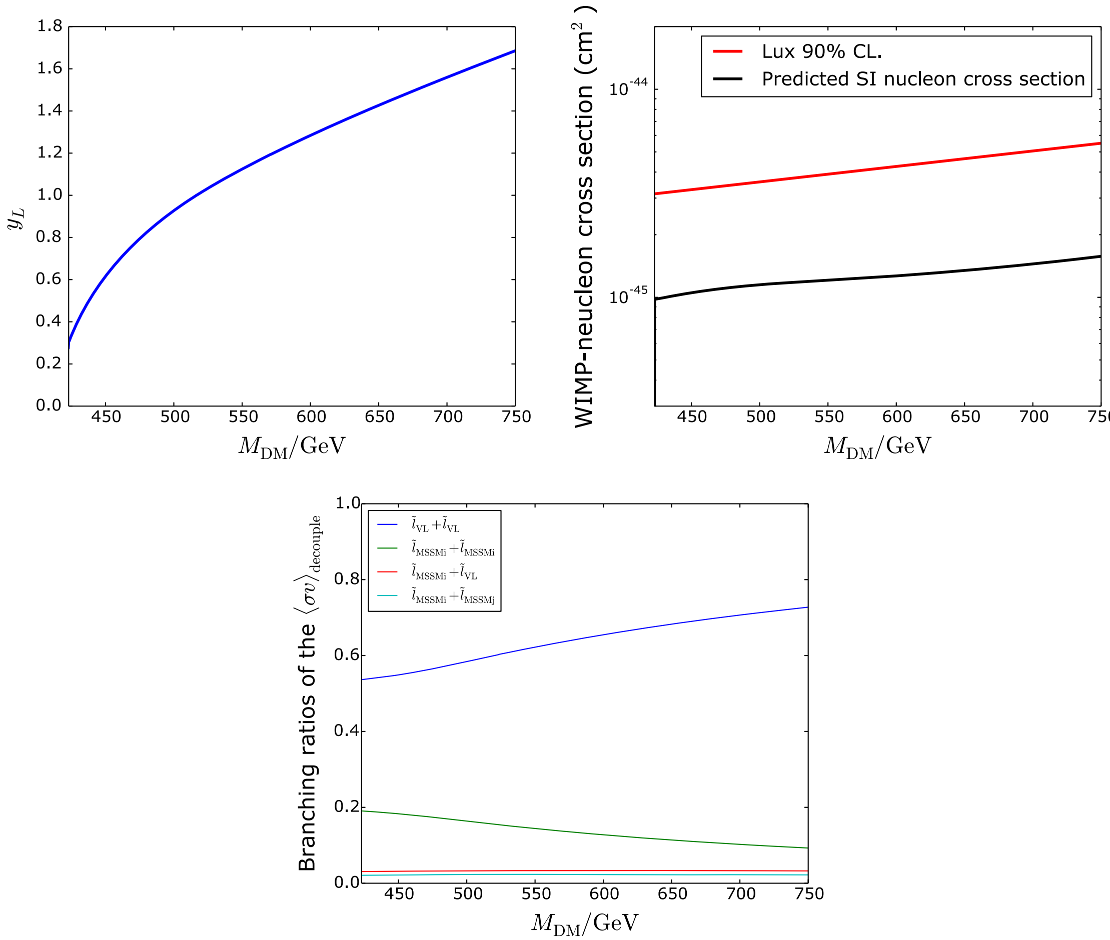

FIG. 3. The y_L corresponding to $\Omega_c h^2 = 0.1199$ (left panel), the spin-independent cross section with the nucleons of the dark matter particles (right panel), and the branching ratios of $\langle\sigma v\rangle_{\rm decouple}$ (bottom panel) in the case in which the vectorlike sleptons coannihilate with two generations of the MSSM slepton.

classifications of the initial states. In Fig. 1, we plot the branching ratios among the coannihilating vectorlike CP-even/CP-odd sneutrino and the vectorlike charged sleptons. In Figs. 2, 3, and 4, we only plot the branching ratios among the vectorlike sleptons and the MSSM sleptons. If we ignore the masses of the MSSM leptons in our numerical calculations, the branching ratios will become generation independent, so we only plot one of the branching ratios of each of the $\tilde{l}_{VL} + \tilde{l}_{\rm MSSMi}$, the $\tilde{l}_{\rm MSSMi} + \tilde{l}_{\rm MSSMi}$, and $\tilde{l}_{\rm MSSMi} + \tilde{l}_{\rm MSSMi} (i \neq j)$ in Figs. 3 and 4.

If the Yukawa coupling constant y_L is switched off, then the main annihilating channels will become the W^+W^-, ZZ channels. The s-channel $R_1 + I_1 \to Z \to \bar{l}l$ is suppressed because the $R - I - Z$ vertex is proportional to $R_1 \partial_\mu I_1 - I_1 \partial_\mu R_1$. At the decoupling time, the four-momentum vector of one dark-side particle is $(m_{\rm DS} + \frac{1}{2} m_{\rm DS} v^2, m_{\rm DS} \vec{v})$. When $v \ll 1$, both terms of $R_1 \partial_\mu I_1 - I_1 \partial_\mu R_1$ nearly cancel out since $m_{R_1} \approx m_{I_1}$.

Generally speaking, if all the coupling constants stay unchanged, the annihilation cross section $\langle\sigma v\rangle_{\rm decouple} \propto \frac{1}{m_{\rm DM}^2}$. If there is only one I_L-like I_1 together with its companions in the same $SU(2)_L$ doublets, that is to say the R_1, the $\tilde{L}^-$, and $\tilde{L}^+$ which coannihilate, $m_{\rm DM} = m_{I_1}$ should be approximately 660 GeV if $y_L \sim 0$. For a heavier m_{I_1}, a larger Yukawa coupling constant y_L is needed for a sufficient $\langle\sigma v\rangle_{\rm decouple} \sim 3 \times 10^{-26}$ cm^3/s. For a lighter m_{I_1}, usually the $\Omega_c h^2$ is suppressed by the too large $\langle\sigma v\rangle_{\rm decouple}$. This can be improved if the MSSM sleptons coannihilate with the vectorlike sleptons. From Fig. 4, we can see that if the dark matter coannihilates with all the MSSM slepton doublets $m_{\rm DM}$ can be as light as ~370 GeV. In the coannihilation scenario, the effective cross section becomes [39]

$$\langle\sigma_{\rm eff} v\rangle = \sum_{ij} \langle\sigma_{ij} v_{ij}\rangle \frac{n_i^{\rm eq}}{n^{\rm eq}} \frac{n_j^{\rm eq}}{n^{\rm eq}}, \tag{22}$$

where i and j indicate the coannihilating particle content. If $\langle\sigma_{ij} v_{ij}\rangle \ll \langle\sigma_{kk} v_{kk}\rangle$ $(i \neq j)$, then $\langle\sigma_{\rm eff} v\rangle$ can be suppressed. In this paper, the cross interactions between the vectorlike sneutrinos and the MSSM sneutrinos can arise from the exchanges of a t-channel Bino or Wino. Thus, heavier

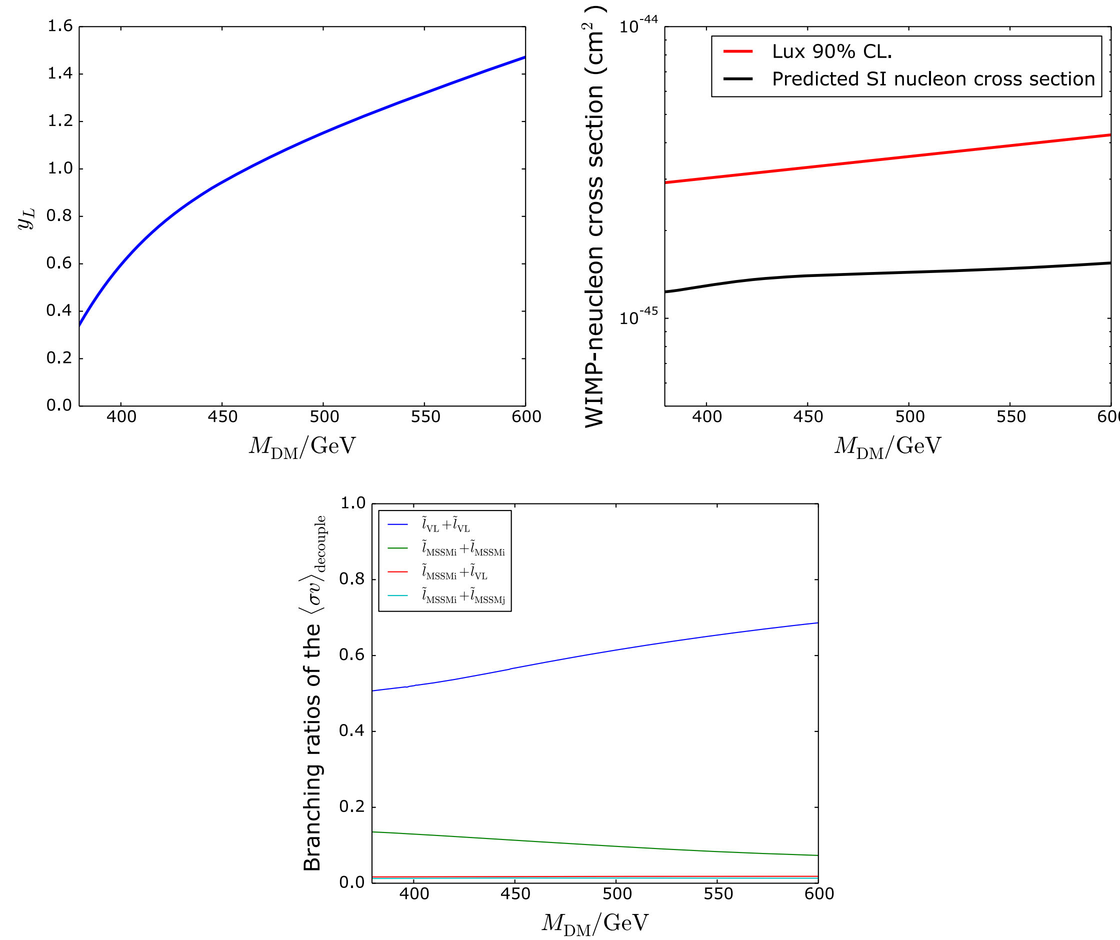

FIG. 4. The y_L corresponding to $\Omega_c h^2 = 0.1199$ (left panel), the spin-independent cross section with the nucleons of the dark matter particles (right panel), and the branching ratios of $\langle \sigma v \rangle_{\rm decouple}$ (bottom panel) in the case in which the vectorlike sleptons coannihilate with all the three generations of the MSSM slepton.

masses of the binos or winos lower the cross interactions and hence lower the $\langle \sigma_{\rm eff} v \rangle$ effectively for the correct relic density in the case of a lighter dark matter. Nevertheless, we should note that the coannihilation scenario requires that $\langle \sigma_{ij} v_{ij} \rangle$ $(i \neq j)$ cannot be too small to avoid the independent annihilation of the different elements; in this case, the masses of the Binos and Winos cannot be too heavy. As has been mentioned before, we adopt the masses of the Binos and Winos to be 2 TeV, which give rise to the cross interactions plotted in Figs. 2, 3, and 4. Further modifying the model can also reach the sufficient cross interactions. For example, in the inverse seesaw model [40–44], the coupling constant y_i in the interaction terms $y_i L_i H_u N$ can be as large as $O(0.1)$, or we can introduce another heavy right-handed neutrino N' as heavy as $\sim 10^{12}$ GeV, and then the coupling constants y'_i, y'_L in the interaction terms $y'_i L_i H_u N'$ and $y'_L L H_u N$ can be as large as $O(0.1)$ (For an example, see the discussions in Appendix B of Ref. [45]). Both these scenarios result in significant $\tilde{L}^\dagger H_u H_u^\dagger \tilde{L}_i$ terms to reach sufficient $\langle \sigma_{L^0 L_i^0} v_{L^0 L_i^0} \rangle$ in order to keep them “co”-annihilating.

As the mass of the dark matter rises up in each coannihilation scenario, the y_L is lifted in order to reach the correct relic density. y_L also contributes to the spin-independent cross section of the dark matter with the nucleons. Various experiments [46–52] have been carried out in order to constrain the dark matter parameters. Among them, we plot the most stringent bound from the LUX [47] in all Figs. 1–4 in comparison with our predicted data. We can see that, although y_L increases as the dark matter mass grows, the constraint line still runs forward the predicted spin-independent cross section.

Finally, we are going to point out that in order to avoid the Landau pole before the gauge coupling constants' unification in a complete $5+\bar{5}$ model y_L should be less than 0.765. This eliminates much area in Figs. 1–4 when the masses of the dark matter particles are heavy. On the other hand, in this situation, the y_L does not make a significant contribution to the SM-like Higgs mass, being unable to relieve the little-hierarchy problem. However, if we relax this condition, the corrections to the SM-like Higgs mass are proportional to y_L^4. If $y_L \sim 1$, and then

$m_L \sim m_{DM}$ is heavy, the Higgs mass can be raised effectively, and we can reach a possible solution to the little-hierarchy problem.

IV. CONCLUSIONS

In place of the MSSM sneutrinos, vectorlike sneutrinos can play the role of dark matter. Compared with the MSSM sneutrinos, the mass splitting between the real part and the imaginary part of the vectorlike sneutrinos can be more naturally acquired without the assumptions of large A terms and do not bother the light neutrino masses. We have calculated the relic density and the elastic scattering cross section with nucleons of the I_L-like dark matter I_1. Coannihilating with the MSSM slepton doublets, the dark matter can be as light as 370 GeV. The predicted cross section with nucleons is also below the most stringent experimental bounds from LUX.

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