

Right-handed sneutrino as cold dark matter

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We consider supersymmetric models with right-handed neutrinos where neutrino masses are *purely Dirac-type*. In this model, right-handed sneutrino can be the lightest supersymmetric particle and can be a viable candidate of cold dark matter of the universe. Right-handed sneutrinos are never thermalized in the early universe because of weakness of Yukawa interaction, but are effectively produced by decays of various superparticles. We show that the present mass density of right-handed sneutrino can be consistent with the observed dark matter density.

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In recent years, various experiments have confirmed the phenomenon of neutrino oscillation. (See, for example, [1–5].) Those results strongly suggest very small but nonvanishing neutrino masses. This fact raises serious problems because the nonvanishing neutrino mass is not allowed in the standard model of particle physics and also because suggested values of neutrino masses are extremely small. The easiest way of generating neutrino masses is to introduce right-handed neutrinos; with this extension, Yukawa couplings of neutrinos may exist. Consequently, neutrinos can acquire masses after electroweak symmetry breaking.

Even with right-handed neutrinos, there are two different classes of scenarios for generating neutrino masses. Probably, more popular one is with Majorana masses for right-handed neutrinos, i.e., so-called “seesaw” scenario [6]. In this scenario, smallness of the neutrino masses is explained by the Majorana masses of right-handed neutrinos which are much larger than the electroweak scale.

Small neutrino masses can be, however, realized without seesaw mechanism. With vanishing Majorana masses of the right-handed neutrinos, which may be the consequence of exact lepton-number symmetry, neutrino masses become Dirac-type. As we will see, in this case, Yukawa coupling constants for neutrinos are, roughly speaking, $O(10^{-13})$ or smaller to make the neutrino masses to be consistent with the results of neutrino-oscillation experiments. One might think that such small Yukawa coupling constants are unnatural. It is, however, natural in 't Hooft's sense [7] since some symmetry (i.e., chiral symmetry in neutrino sector) is restored in the limit of vanishing neutrino Yukawa coupling constants. In this model, however, right-handed neutrinos are mostly irrelevant for collider experiments and cosmology since their interaction is extremely weak.

In supersymmetric models, which is strongly motivated as a solution to various problems in the standard model of particle physics, like hierarchy and naturalness problems, situation changes. In particular, superpartners of right-handed neutrinos may play an important role in cosmology. When small neutrino masses are purely Dirac-type, masses

of right-handed sneutrinos are dominantly from effects of supersymmetry (SUSY) breaking. Then, one should note that the lightest superparticle (LSP) may be the (lightest) right-handed neutrino $\tilde{\nu}_R$. Importantly, since the LSP $\tilde{\nu}_R$ becomes stable by the R -parity conservation and also is very weakly interacting, it can be a viable candidate of cold dark matter (CDM) provided that its relic density is the right amount.

In this paper, we consider the minimal supersymmetric standard model (MSSM) with three generations of right-handed (s)neutrinos where small neutrino masses are purely Dirac-type. In particular, we study the case where the LSP is the lightest right-handed sneutrino and see if the relic density of $\tilde{\nu}_R$ can become consistent with the present CDM density. Since interaction of right-handed sneutrino is very weak, it is not thermalized in thermal bath. Even in this case, some of decay (and scattering) processes produce $\tilde{\nu}_R$, as in the cases of other SuperWIMPs. As we will see, in some parameter region, the density parameter of right-handed sneutrino $\Omega_{\tilde{\nu}_R}$ can be $O(0.1)$, which is consistent with the CDM density suggested by the WMAP [8]:

$$\Omega_c h_{100}^2 = 0.1126_{-0.0181}^{+0.0161}, \quad (1)$$

where h_{100} is the Hubble constant in units of 100km/sec/Mpc.

Let us first introduce interaction and mass terms in Lagrangian. The important part of superpotential is

$$W = y_\nu \hat{H}_u \hat{L} \hat{\nu}_R^c + \mu_H \hat{H}_u \hat{H}_d, \quad (2)$$

where $\hat{H}_u = (\hat{H}_u^+, \hat{H}_u^0)$ and $\hat{H}_d = (\hat{H}_d^0, \hat{H}_d^-)$ are up- and down-type Higgses, and $\hat{L} = (\hat{\nu}_L, \hat{l}_L^-)$ the left-handed lepton. (In this letter, “hat” is for superfields while “tilde” is for superparticles with odd R -parity.) Here, we omit flavor indices for simplicity. With this superpotential, neutrino mass is generated after electroweak symmetry breaking: $m_\nu = y_\nu \langle H_u^0 \rangle = y_\nu v \sin\beta$, where $v \simeq 174$ GeV and $\tan\beta = \langle H_u^0 \rangle / \langle H_d^0 \rangle$. Thus, the Yukawa coupling constant y_ν is determined once the neutrino mass is fixed:

$$y_\nu \sin\beta = 3.0 \times 10^{-13} \times \left(\frac{m_\nu^2}{2.8 \times 10^{-3} \text{ eV}^2} \right)^{1/2}. \quad (3)$$

For simplicity, we consider the case where $\tan\beta$ is relatively large. In this case H_u behaves like the standard-model Higgs; $H_u \simeq H_{SM}$. The lightest Higgs boson h is contained in H_{SM} , and we take its mass to be $m_h = 115$ GeV in our study. On the contrary, H_d plays the role of heavy Higgs doublet. In addition, we approximate mass eigenstates of the superparticles by left- and right-handed sleptons, gauginos, and Higgsinos. Effects of left-right mixing are taken into account by mass-insertion method. The mass terms for superparticles are given by

$$\mathcal{L}_{\text{mass}} = -\frac{1}{2}[m_{\tilde{B}}\tilde{B}\tilde{B} + m_{\tilde{W}}\tilde{W}\tilde{W} + \mu_H\tilde{H}_u\tilde{H}_d + \text{h.c.}] - m_{\tilde{\nu}_R}^2\tilde{\nu}_R^*\tilde{\nu}_R - m_{\tilde{l}_L}^2\tilde{l}_L^*\tilde{l}_L - m_{\tilde{\nu}_L}^2\tilde{\nu}_L^*\tilde{\nu}_L, \quad (4)$$

where \tilde{B} and \tilde{W} are gauginos for $U(1)_Y$ and $SU(2)_L$ gauge groups, respectively. Notice that all the terms except for Higgsino mass term are from SUSY breaking and that $m_{\tilde{l}_L}^2$ and $m_{\tilde{\nu}_L}^2$ are identical up to D -term contribution. Furthermore, we introduce soft-SUSY breaking trilinear coupling which will become important in the following discussion

$$\mathcal{L}_A = A_\nu H_u \tilde{L} \tilde{\nu}_R^c + \text{h.c.} \quad (5)$$

For later convenience, we parameterize

$$m_{\tilde{l}_L, \tilde{\nu}_L}^2 = m_{\tilde{L}}^2 + \Delta_{\tilde{l}_L, \tilde{\nu}_L}^{(D)}, \quad A_\nu = a_\nu y_\nu m_{\tilde{L}}, \quad (6)$$

where $\Delta_{\tilde{l}_L, \tilde{\nu}_L}^{(D)}$ represents the D -term contribution. In addition, we adopt the GUT relation among the gaugino masses: $m_{\tilde{B}} \simeq 0.5m_{\tilde{W}}$.

Now let us discuss the production of the LSP right-handed sneutrino $\tilde{\nu}_R$ via various decay processes. Since we found that the production via scatterings is subdominant and can be neglected, we consider here only the decay processes $x \rightarrow \tilde{\nu}_R y$. We assume that particles x and y are in chemical equilibrium. Their distribution functions at the temperature T , then, are given by $f(E) = (e^{E/T} \pm 1)^{-1}$, where the positive and negative signs are for fermions and bosons, respectively.

The evolution of the number density of $\tilde{\nu}_R$ is governed by the Boltzmann equation

$$\dot{n}_{\tilde{\nu}_R} + 3Hn_{\tilde{\nu}_R} = C_{\text{decay}}, \quad (7)$$

where ‘‘dot’’ represents time-derivative and H is the Hubble parameter. The source term from the decay processes is written in the form

$$C_{\text{decay}} = \sum_{x,y} \int \frac{d^3k_x}{(2\pi)^3} \gamma_x (2s_x + 1) \Gamma_{x \rightarrow \tilde{\nu}_R y} f_x \langle 1 \pm f_y \rangle_{k_x}, \quad (8)$$

where $\gamma_x = m_x / \sqrt{k_x^2 + m_x^2}$ is the Lorentz factor and $(2s_x + 1)$ is the spin multiplicity of x . Furthermore, $\langle 1 \pm f_y \rangle_{k_x}$ is the averaged final-state multiplicity factor for fixed value

of initial-state momentum. (Here, the positive and negative signs are for bosons and fermions, respectively.)

Decay rates for processes which become important for the calculation of Ω_{ν_R} are given by

$$\Gamma_{\tilde{H}^0 \rightarrow \tilde{\nu}_R \tilde{\nu}_L} = \Gamma_{\tilde{H}^+ \rightarrow \tilde{\nu}_R \tilde{l}_L^+} = \frac{\beta_f^2 y_\nu^2}{32\pi} \mu_H, \quad (9)$$

$$\Gamma_{\tilde{\nu}_L \rightarrow \tilde{\nu}_R h} = \frac{\beta_f}{32\pi} \frac{A_\nu^2}{m_{\tilde{\nu}_L}}, \quad (10)$$

$$\Gamma_{\tilde{\nu}_L \rightarrow \tilde{\nu}_R Z} = \frac{\beta_f^3}{32\pi} \left[\frac{m_{\tilde{\nu}_L}^2}{m_{\tilde{\nu}_L}^2 - m_{\tilde{\nu}_R}^2} \right]^2 \frac{A_\nu^2}{m_{\tilde{\nu}_L}}, \quad (11)$$

$$\Gamma_{\tilde{l}_L \rightarrow \tilde{\nu}_R W^-} = \frac{\beta_f^3}{16\pi} \left[\frac{m_{\tilde{l}_L}^2}{m_{\tilde{\nu}_L}^2 - m_{\tilde{\nu}_R}^2} \right]^2 \frac{A_\nu^2}{m_{\tilde{l}_L}}, \quad (12)$$

$$\Gamma_{\tilde{B} \rightarrow \tilde{\nu}_R \tilde{\nu}_L} = \frac{\beta_f^2 g_1^2}{64\pi} \left[\frac{A_\nu v_T}{m_{\tilde{\nu}_L}^2 - m_{\tilde{\nu}_R}^2} \right]^2 m_{\tilde{B}}, \quad (13)$$

$$\Gamma_{\tilde{W}^0 \rightarrow \tilde{\nu}_R \tilde{\nu}_L} = \frac{\beta_f^2 g_2^2}{64\pi} \left[\frac{A_\nu v_T}{m_{\tilde{\nu}_L}^2 - m_{\tilde{\nu}_R}^2} \right]^2 m_{\tilde{W}}, \quad (14)$$

$$\Gamma_{\tilde{W}^+ \rightarrow \tilde{\nu}_R \tilde{l}_L^+} = \frac{\beta_f^2 g_2^2}{32\pi} \left[\frac{A_\nu v_T}{m_{\tilde{\nu}_L}^2 - m_{\tilde{\nu}_R}^2} \right]^2 m_{\tilde{W}}, \quad (15)$$

where we have neglected lepton masses. Here, g_1 and g_2 are gauge coupling constants for the $U(1)_Y$ and $SU(2)_L$ gauge groups, respectively, and, for the process $x \rightarrow \tilde{\nu}_R y$, β_f is given by

$$\beta_f^2 = \frac{1}{m_x^4} [m_x^4 - 2(m_{\tilde{\nu}_R}^2 + m_y^2)m_x^2 + (m_{\tilde{\nu}_R}^2 - m_y^2)^2], \quad (16)$$

with m_x and m_y being the masses of the particles x and y , respectively. In addition, v_T is temperature-dependent Higgs VEV. We approximate the Higgs potential in thermal bath as [9]

$$V_T \simeq \frac{m_h^2}{4v^2} (|H_{SM}|^2 - v^2)^2 + \frac{1}{8v^2} (2m_{\tilde{W}}^2 + m_Z^2 + 2m_t^2) T^2 |H_{SM}|^2, \quad (17)$$

where T is cosmic temperature. In the calculation of V_T , we take into account only the effects of standard-model particles since the production of $\tilde{\nu}_R$ becomes effective when the cosmic temperature is lower than the masses of relevant superparticles. Then we minimize V_T to obtain $v_T \simeq \langle H_u \rangle_T$.

In solving Eq. (7) it is useful to define the yield variable

$$Y_{\tilde{\nu}_R} \equiv \frac{n_{\tilde{\nu}_R}}{s}, \quad (18)$$

where $s \equiv \frac{2\pi^2}{45} g_* T^3$ is the total entropy density of the universe. (In this analysis, we take the effective number of massless degrees of freedom as $g_* = 106.75$, since the production of $\tilde{\nu}_R$ becomes effective at the temperature lower than most of superparticle masses.)

Using the relation $\dot{T} = -HT$, $Y_{\tilde{\nu}_R}$ is given by

$$Y_{\tilde{\nu}_R}(T) = \int_T^{T_{\max}} \frac{C_{\text{decay}}}{sHT} dT. \quad (19)$$

Once the yield variable is obtained, we can also calculate the mass density of sneutrino using the relation $\rho_{\tilde{\nu}_R}/s = m_{\tilde{\nu}_R} Y_{\tilde{\nu}_R}$. Comparing with the present value of the critical density ρ_{crit} , which is given by $[\rho_{\text{crit}}/s]_{\text{now}} \simeq 3.6h_{100}^2 \times 10^{-9}$ GeV, we obtain density parameter of right-handed sneutrino $\Omega_{\tilde{\nu}_R} \equiv (\rho_{\tilde{\nu}_R} + \rho_{\tilde{\nu}_R^*})/\rho_{\text{crit}}$.

Importantly the present value of $Y_{\tilde{\nu}_R}$ is insensitive to the maximal temperature T_{\max} of the universe (as long as T_{\max} is much higher than the masses of superparticles). To see this, it is instructive to roughly estimate $Y_{\tilde{\nu}_R}$. Neglecting the Lorentz factor, let us approximate the decay term as $C_{\text{decay}} \sim \frac{N_{\text{mode}}}{16\pi} y_\nu^2 m_{\text{SUSY}} n_{\text{SUSY}}$, where m_{SUSY} and n_{SUSY} are typical mass scale and number density of (parent) superparticles, respectively, and N_{mode} is the number of possible decay mode. The number density n_{SUSY} is $\sim T^3$ when the temperature is higher than m_{SUSY} , and is exponentially suppressed when $T \ll m_{\text{SUSY}}$. The integration in Eq. (19) is then dominated at the temperature of $T \sim m_{\text{SUSY}}$. Consequently, we obtain $Y_{\tilde{\nu}_R}(T \ll m_{\text{SUSY}}) \sim \frac{N_{\text{mode}} y_\nu^2 M_*}{16\pi g_*^{3/2} m_{\text{SUSY}}}$, with $M_* \simeq 2.4 \times 10^{18}$ GeV being the reduced Planck scale. As one can see, the present $Y_{\tilde{\nu}_R}$ is insensitive to thermal history for $T \gg m_{\text{SUSY}}$. Thus, $\Omega_{\tilde{\nu}_R}$ does not depend on, for example, reheating temperature after inflation. We can also estimate the density parameter; using the above estimation of $Y_{\tilde{\nu}_R}$, we obtain $\Omega_{\tilde{\nu}_R} h_{100}^2 \sim O(10^{-3}) \times N_{\text{mode}} (\frac{y_\nu}{3 \times 10^{-13}})^2 (\frac{m_{\tilde{\nu}_R}}{m_{\text{SUSY}}})$. With this naive estimation, $\Omega_{\tilde{\nu}_R}$ becomes smaller than the currently observed CDM density. In some case, however, $\Omega_{\tilde{\nu}_R}$ becomes much larger, as we see below.

Now, we are at the position to quantitatively estimate $\Omega_{\tilde{\nu}_R}$. We have numerically evaluated the yield variable using Eq. (19) for several choices of parameters. The relic density of right-handed sneutrino strongly depends on the neutrino Yukawa coupling constant which is related to the neutrino mass. Neutrino-oscillation experiments determine only the mass-squared differences of neutrinos; here we adopt [3,5]

$$[\Delta m_\nu^2]_{\text{atm}} = 2.8 \times 10^{-3} \text{ eV}^2, \quad (20)$$

$$[\Delta m_\nu^2]_{\text{solar}} = 7.9 \times 10^{-5} \text{ eV}^2. \quad (21)$$

First, let us consider the case where the masses of the neutrinos are hierarchical. In this case, $[\Delta m_\nu^2]_{\text{atm}}$ almost

corresponds to the mass-squared of the heaviest neutrino which we call third-generation neutrino. In this case, the largest Yukawa coupling constant is given by $y_\nu^{(3)} \simeq 3 \times 10^{-13}$, and other Yukawa coupling constants are much smaller. (Here and hereafter, the superscript “(i)” indicates that $y_\nu^{(i)}$ is for neutrino in i -th generation.) As a result, production of $\tilde{\nu}_R$ is dominated by processes where the third-generation (s)neutrino is related.

As show in in Eqs. (9)–(15), there exist various decay processes which produce $\tilde{\nu}_R$. Among them, Higgsino decay process dominates the $\tilde{\nu}_R$ production when the effects of the trilinear scalar couplings are negligible. In this case, however, $\Omega_{\tilde{\nu}_R}$ becomes too small to be consistent with observation of the CDM relic density. Indeed, when $m_{\tilde{\nu}_R} = 100$ GeV, for example, we found $\Omega_{\tilde{\nu}_R} = 0.004 - 0.001$ for $\mu_H = 200$ GeV – 1 TeV, which is much smaller than the present CDM density.

If some of other decay processes are effective, $\Omega_{\tilde{\nu}_R}$ can become significantly larger. In particular, when the trilinear coupling constant A_ν is nonvanishing, decays of various superparticles produces $\tilde{\nu}_R$ sufficiently.

One possible enhancement can be due to the process $\tilde{W} \rightarrow \tilde{\nu}_R + \dots$; if the mass of left-handed sneutrino $\tilde{\nu}_L$ is close to $m_{\tilde{\nu}_R}$, production of $\tilde{\nu}_R$ is enhanced because the left-right mixing of sneutrinos becomes larger as $|m_{\tilde{\nu}_L}^2 - m_{\tilde{\nu}_R}^2|$ becomes smaller. (See Eqs. (13)–(15)) In fact, $\Omega_{\tilde{\nu}_R}$ can be of $O(0.1)$ with a mild degeneracy of $\tilde{\nu}_R$ and $\tilde{\nu}_L$. To see this, in Fig. 1, we show $\Omega_{\tilde{\nu}_R} h_{100}^2$ as a function of $m_{\tilde{\nu}_L}$. We can see that 10 – 20% degeneracy is enough to realize $\tilde{\nu}_R$ -CDM even if $a_\nu \leq 3$. Here, we take $m_{\tilde{W}} = 300$ GeV. If we increase the Wino mass, $\Omega_{\tilde{\nu}_R}$ decreases since the

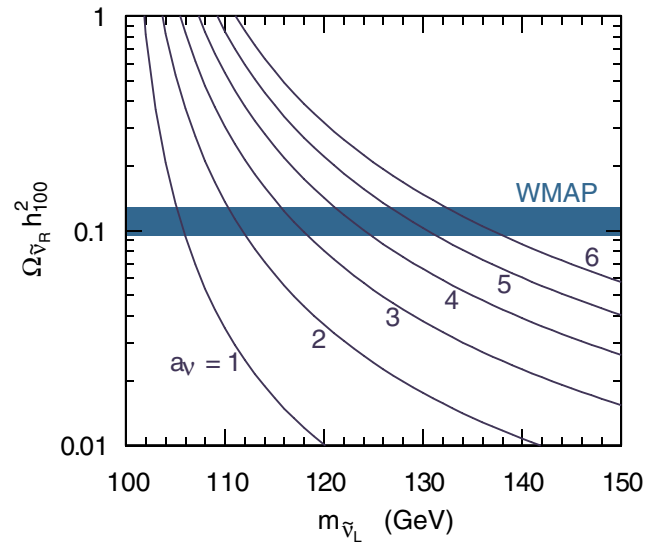


FIG. 1 (color online). $\Omega_{\tilde{\nu}_R} h_{100}^2$ as a function of $m_{\tilde{\nu}_L}$ for $a_\nu = 1 - 6$. Here, we take $m_{\tilde{\nu}_R} = 100$ GeV, $m_{\tilde{W}} = 300$ GeV, and $\mu_H = 150$ GeV. The shaded band corresponds to the CDM density suggested by the WMAP.

temperature-dependent VEV v_T becomes suppressed at high temperature. We also note that the process $\tilde{H}^+ \rightarrow \tilde{\nu}_R \tau_R^+$ gives additional contribution to $\Omega_{\tilde{\nu}_R}$, which may become sizable when $\tan\beta$ is large. This is because the partial width of this decay process is proportional to $\tan^2\beta$. We found that $\Omega_{\tilde{\nu}_R}$ can be $\sim 40\%$ larger for $\tan\beta = 55$.

Even without mass degeneracy between $\tilde{\nu}_R$ and $\tilde{\nu}_L$, $\tilde{\nu}_R$ -CDM is realized if $a_\nu \gg 1$. Although such a very large value of a_ν may not be realized in simple supergravity models, it is phenomenologically viable.

Next, we turn to consider the case of degenerate neutrino masses, where there is another possibility of enhancing $\Omega_{\tilde{\nu}_R}$. In this case, the relation $y_\nu^{(1)} \simeq y_\nu^{(2)} \simeq y_\nu^{(3)}$ holds and all three generations of right-handed sneutrinos may be effectively produced. The point is that the neutrino Yukawa coupling constants can be much larger than $\sim 3 \times 10^{-13}$, and hence $\Omega_{\tilde{\nu}_R}$ can be more enhanced than the hierarchical case. $\tilde{\nu}_R$ -CDM can be realized if the neutrino masses are $O(0.1 \text{ eV})$ even if there is no other enhancement.

So far, we have only discussed the production of $\tilde{\nu}_R$ by the decay of superparticles *in chemical equilibrium*. In general, $\tilde{\nu}_R$ can be also produced after superparticles freeze out. The lightest superparticle in the MSSM sector, which we call MSSM-LSP, is very long-lived because of the weakness of the neutrino Yukawa interaction. Thus, decay of the MSSM-LSP occurs after its number is frozen, which gives another source of $\tilde{\nu}_R$. The total density parameter is then given by $\Omega_{\tilde{\nu}_R}^{(\text{total})} = \Omega_{\tilde{\nu}_R} + \frac{m_{\tilde{\nu}_R}}{m_{\text{MSSM}}} \Omega_{\text{MSSM}}$, where m_{MSSM} is the MSSM-LSP mass. Here, the first term is evaluated with Eq. (19) while the second one is from the decay after the freeze out; Ω_{MSSM} is the expected density parameter of the MSSM-LSP for the case where the MSSM-LSP is stable. Notice that Ω_{MSSM} is model-dependent and may be smaller than the WMAP value; if Ω_{MSSM} is large, the parameter space to realize $\tilde{\nu}_R$ -CDM is enlarged.

Here we discuss the effects of the MSSM-LSP decay on big-bang nucleosynthesis (BBN). If the MSSM-LSP decays after BBN starts, it may cause hadro- and photo-dissociation of light elements, which may spoil the success of standard BBN scenario. In order to realize $\tilde{\nu}_R$ -CDM with enhanced left-right mixing, $\tilde{\nu}_L$ is required to be relatively light and then it may be the MSSM-LSP. If so, we found that, for $m_{\tilde{\nu}_R} = 100 \text{ GeV}$, its lifetime becomes so long that light-element abundances become inconsistent with observations due to hadro-dissociation processes [10] unless $m_{\tilde{\nu}_L} \gtrsim 140 \text{ GeV}$ (with $a_\nu \gtrsim 6$). On the other hand, very degenerate sneutrinos ($m_{\tilde{\nu}_L} - m_{\tilde{\nu}_R} \lesssim \text{a few GeV}$) is another possibility to avoid the difficulty by suppressing the energy release from $\tilde{\nu}_L$ decay. Otherwise, the BBN constraints can be also avoided if the MSSM-LSP is the lightest neutralino, or if R -parity in the MSSM sector is (very weakly) broken.

Finally, we comment on the decay of heavier right-handed sneutrinos (denoted by $\tilde{\nu}'_R$). When the mass differ-

ence between $\tilde{\nu}'_R$ and $\tilde{\nu}_R$ is large enough, $\tilde{\nu}'_R$ may decay, for example, into $\tilde{H}_u L'$, and the produced Higgsino decays subsequently as $\tilde{H}_u \rightarrow \tilde{\nu}_R L$. In this case, lifetime of the long-lived particles can be $\sim 1 \text{ sec}$ or shorter; such decay processes are cosmologically safe [10]. On the contrary, if $\tilde{\nu}'_R$ decays only through the processes of $\tilde{\nu}'_R \rightarrow \tilde{\nu}_R + \dots$ due to kinematical reason, the lifetime becomes much longer than the present age of the universe. For instance, when $m_{\tilde{\nu}_R} = 100 \text{ GeV}$, $m_{\tilde{\nu}'_R} = 150 \text{ GeV}$, and $\mu_H = 200, 300, \text{ and } 400 \text{ GeV}$, the lifetime of $\tilde{\nu}'_R$ is estimated as $\tau_{\tilde{\nu}'_R} \simeq 5 \times 10^{24}, 5 \times 10^{25}, \text{ and } 2 \times 10^{26} \text{ yr}$, respectively. (Here, we approximated that all the neutrino Yukawa coupling constants are 10^{-13} , and neglected trilinear scalar couplings.) Even in this case, some (tiny) amount of heavier right-handed sneutrinos decay into charged particles until today. Assuming the density parameter of $\tilde{\nu}'_R$ to be $O(0.01)$, flux of the (primary) charged particles is estimated to be $\mathcal{F} \sim 10^{-16} \text{ cm}^{-2} \text{ sec}^{-1} \text{ str}^{-1} \text{ GeV}^{-1} \times \left(\frac{\tau_{\tilde{\nu}'_R}}{10^{25} \text{ yr}}\right)^{-1}$. The emitted charged particles scatter off the background radiation and produce energetic photons. However, we expect that such small flux is consistent with the observations since the above flux is much smaller than the observed γ -ray flux with the energy $10 - 100 \text{ GeV}$: $\mathcal{F}^{(\text{obs})} \sim 10^{-8} - 10^{-10} \text{ cm}^{-2} \text{ sec}^{-1} \text{ str}^{-1} \text{ GeV}^{-1}$ [11].

In summary, in this letter, we have considered supersymmetric models with the right-handed-sneutrino LSP, where masses of neutrinos are purely Dirac-type. We have seen that, in such models, $\Omega_{\tilde{\nu}_R}$ can be of $O(0.1)$ and right-handed sneutrino can be CDM. In particular, if $\tilde{\nu}_R$ -CDM is realized in the parameter region where $\tilde{\nu}_R$ and $\tilde{\nu}_L$ are quite degenerate, left-handed sleptons become fairly light. Such light sleptons are interesting targets of the future collider experiments. On the contrary, if $\tilde{\nu}_R$ -CDM is realized with degenerate neutrino masses, the (active) neutrino mass should become $\sim O(0.1 \text{ eV})$. Such a scenario may be tested by using matter power spectrum from Ly- α forest data [12].

If $\tilde{\nu}_R$ is the LSP, there are much more to be discussed from phenomenological and cosmological points of view. Even though the LSP is $\tilde{\nu}_R$, the MSSM-LSP looks like stable particle in collider experiments. However, in this case, the MSSM-LSP may be charged or even colored. In addition, cosmology may be also significantly affected if the right-handed sneutrino is the LSP; one of the examples has been discussed in this letter. We also note that, if the initial amplitude of $\tilde{\nu}_R$ is $\sim 10^9 \text{ GeV}$, coherent mode of the right-handed sneutrino may become the CDM. We believe that the scenario with right-handed-sneutrino LSP provides new and rich phenomenology.

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