Cosmic rays from Leptonic Dark Matter

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Abstract

If dark matter possesses a lepton number, it is natural to expect the dark-matter annihilation and/or decay mainly produces the standard model leptons, while negligible amount of the antiproton is produced. To illustrate such a simple idea, we consider a scenario that a right-handed sneutrino dark matter decays into the standard model particles through tiny $R$-parity violating interactions. Interestingly enough, charged leptons as well as neutrinos are directly produced, and they can lead to a sharp peak in the predicted positron fraction. Moreover, the decay of the right-handed sneutrino also generates diffuse continuum gamma rays which may account for the excess observed by EGRET, while the primary antiproton flux can be suppressed. Those predictions on the cosmic-ray fluxes of the positrons, gamma rays and antiprotons will be tested by the PAMELA and FGST observatories.
I. INTRODUCTION

The presence of dark matter has been securely established by numerous observational evidences. In particular, the latest 5yr WMAP data determined the dark matter abundance with unprecedented precision as

\[ \Omega_{DM}h^2 = 0.1099 \pm 0.0062, \]  

where \( \Omega_{DM} \) is the fraction of critical density in dark matter and \( h \) is the Hubble parameter in units of 100 km/Mpc/sec. In the meantime it is not known yet what dark matter is made of despite many experimental direct/indirect searches hitherto — but there is a hope that the PAMELA [2] and FGST (formerly GLAST) [3] satellites may provide us with some important information on the nature of dark matter.

It is normally assumed that the dark matter is charged under an exact discrete symmetry to ensure its stability, for instance, \( R \)-parity in a supersymmetric theory. There might exist many discrete symmetries realized in our vacuum and one of them may be responsible for making the dark matter stable. However, if the discrete symmetry is broken, the dark matter will be unstable and eventually decay into the Standard Model (SM) particles. Recently such a decaying dark matter has attracted much attention (e.g. the gravitino with \( R \)-parity violation [4, 5, 6, 7, 8] or a hidden \( U(1) \) gauge boson [9] and gaugino [10]), since the high-energy cosmic rays produced by the decay of dark matter may account for the observed anomalous excesses in gamma rays [11, 12] and/or positrons [13].

Very recently, the PAMELA data on the antiproton flux has been released [14], and it suggests that the observed antiprotons are mainly secondaries, which is consistent with the balloon-borne experiments at the top of atmosphere (TOA) [15]. This will place tight constraints on possible dark matter candidates for explaining the anomalous excesses in positrons and/or gamma rays [16]. Of course, the predicted antiproton flux has a large uncertainty mainly due to our poor understanding of the cosmic-ray propagation inside our Galaxy. However, if we take the constraint on the antiproton flux seriously, and if we attribute the excesses in positrons and gamma rays to the dark-matter annihilation or decay, we are led to explore a dark matter candidate naturally satisfying the constraint on the antiproton flux. Our main idea is as follows. If the dark matter has a lepton number, it is quite natural to expect that it annihilates or decays mainly into the leptons, while only
negligible amount of the antiprotons are produced. The purpose of this paper is to illustrate this simple and naive idea by using an explicit example.

There are several candidates for leptonic dark matter such as the sterile neutrino [17]. As an example, in this paper we consider a scenario that a right-handed sneutrino $\tilde{\nu}_R$ is the lightest supersymmetric particle (LSP) and accounts for dark matter of the universe. The neutrino mass is assumed to be Dirac type in our study. If the $R$-parity is an exact symmetry of nature, the $\tilde{\nu}_R$ dark matter is absolutely stable [18, 19, 20]. However, the $R$-parity may not be an exact symmetry and is only an approximate one accompanied with tiny violations. We focus on the $R$-parity violating bilinear term throughout this paper. In the absence of the $R$-parity, the $\tilde{\nu}_R$ dark matter is not stable anymore, and directly decays into the charged leptons ($\tau$, $\mu$ and $e$) and neutrinos, and it can also decay into the quarks and the $W$ and $Z$ gauge bosons depending on the mixing with the Higgs bosons. As a result, a sharp peak is predicted in the positron fraction, which may explain the anomalous excess observed by High Energy Antimatter Telescope (HEAT) [13], MASS [21] and AMS [22] experiments. Interestingly, the preliminary PAMELA data also exhibits such anomalous excess in the cosmic-ray positron fraction, although we do not try to include the preliminary data since it is not formally released yet. The continuum gamma rays are also produced mainly from the decay of pions, and those gamma rays may account for the excesses observed by EGRET [11]. For a proper set of parameters, the production of antiprotons can be suppressed, which is consistent with an observational fact that the antiprotons measured by the balloon-borne experiments [15] and also by the PAMELA satellite [14] are considered to be mainly secondaries [23]. The suppression in the antiproton flux is particularly important because some decaying dark matter scenarios predict too large antiproton flux at the solar system [7].

This paper is organized as follows. In Sec. II, we present the effective interactions and the relevant decay modes used in our calculations. In Sec. III, we calculate the spectra of the positron, gamma-ray, and antiproton produced from the decay of the $\tilde{\nu}_R$, and compare them with the observational data. We also discuss how the right-handed sneutrinos are produced in the early universe in Sec. IV, and we give our discussions and conclusions in Sec. V.
II. FRAMEWORK

The non-vanishing neutrino masses have been firmly established by neutrino oscillation experiments (see Ref. [24] for recent review and references therein), although it is not known yet whether the neutrinos are Majorana or Dirac fermions. In this paper, we consider a Dirac-type neutrino mass in a supersymmetry theory without $R$-parity conservation, and assume the gravity mediation. We introduce two additional interactions to the minimal supersymmetric standard model (MSSM); one is the neutrino Yukawa coupling for the Dirac neutrino mass, and the other is the $R$-parity violating bilinear term. The superpotential is therefore

$$W = W_{\text{MSSM}} + y^\nu_{ij} \bar{N}_i L_j H_u + \mu_i H_u L_i,$$

(2)

$$W_{\text{MSSM}} = y^u_i \bar{U}_i Q_i H_u - y^d_i \bar{D}_i Q_i H_d - y^e_i \bar{E}_i L_i H_d + \mu H_u H_d,$$

(3)

where $\bar{N}$ is the right-handed neutrino superfield, $y^\nu$ is the neutrino Yukawa coupling, $\mu_i$ denotes the coefficient of the $R$-parity violating bilinear term, and the indices $i$ and $j$ denote the generations. We neglect the flavor mixings in the MSSM Yukawa interactions, and assume the minimal Kähler potential for all the MSSM particles as well as the right-handed neutrinos.

Let us first discuss the neutrino Yukawa coupling. The three left-handed neutrinos in the weak eigenstate $\nu_i$ are related with the mass eigenstates $\tilde{\nu}_i$ as

$$\nu_i = U_{ij} \cdot \tilde{\nu}_j,$$

(4)

where $i$ runs from 1 to 3 and the well-known mixing matrix $U$ takes the following form:

$$U = \begin{pmatrix}
    c_{13}c_{12} & s_{12}c_{13} & s_{13} \\
    -s_{12}c_{23} - s_{23}s_{13}c_{12} & c_{23}c_{12} - s_{23}s_{13}s_{12} & s_{23}c_{13} \\
    s_{23}s_{12} - s_{13}c_{23}c_{12} & -s_{23}c_{12} - s_{13}s_{12}c_{23} & c_{23}c_{13}
\end{pmatrix},$$

(5)

where $c_{ij} \equiv \cos \theta_{ij}$ and $s_{ij} \equiv \sin \theta_{ij}$, and $\theta_{ij}$ denotes the mixing angle of neutrinos $\nu_i$ and $\nu_j$.

We can define the right-handed neutrinos so that the neutrino Yukawa coupling matrix is

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#1 The right-handed sneutrinos acquire tiny vacuum expectation values (vev) after the electroweak symmetry breaking in the presence of the bilinear $R$-parity violating term. However, the expectation values are suppressed by the neutrino Yukawa couplings, and so, we can neglect their effects. On the other hand, the bilinear term can be induced from the neutrino Yukawa couplings with non-vanishing vevs of the right-handed sneutrinos or the linear terms in the Kähler potential.
TABLE I: The best-fit and adopted values of three-flavor neutrino oscillation parameters from global data, including solar, atmospheric, reactor (KamLAND, CHOOZ) and accelerator (K2K) experiments [25].

|                | $\Delta m^2_{21}$ [eV$^2$] | $|\Delta m^2_{31}|$ [eV$^2$] | $\sin^2 \theta_{12}$ | $\sin^2 \theta_{23}$ | $\sin^2 \theta_{13}$ |
|----------------|-----------------------------|-----------------------------|------------------------|------------------------|------------------------|
| best fit       | $7.65^{+0.23}_{-0.20} \times 10^{-5}$ | $2.40^{+0.12}_{-0.11} \times 10^{-3}$ | 0.304$^{+0.022}_{-0.016}$ | 0.50$^{+0.07}_{-0.06}$ | 0.01$^{+0.016}_{-0.011}$ |
| adopted        | $7.65 \times 10^{-5}$        | $2.4 \times 10^{-3}$        | 0.304                  | 0.50                  | 0                      |

given by

$$y^\nu = \text{diag}(m_1, m_2, m_3) U^\dagger,$$

where $m_i$ is the mass of $\hat{\nu}_i$. In Table I, we show the observational constraints on those mixings and the mass differences based on the global three-neutrino analysis [25]. Since the neutrino oscillation data is not sensitive to the absolute masses of the neutrinos, the following three neutrino mass spectra are possible: (i) normal hierarchy case $m_3 > m_2 \gg m_1$ (ii) inverted hierarchy case $m_2 > m_1 > m_3$ and (iii) degenerate case $m_1 \simeq m_2 \simeq m_3$. Throughout this paper, we adopt the normal hierarchy for simplicity, therefore, $m_2^2 \sim \Delta m^2_{21}$, $m_3^2 \sim \Delta m^2_{31}$ and we adopt massless $m_1$ in our numerical study.

Next we consider the $R$-parity violation. In the presence of the bilinear $R$-parity violation, it is known that there is no unique way to define $L_i$ and $H_d$ since they have the same quantum numbers [26]. Taking account of the soft SUSY breaking terms, the left-handed sneutrino $\langle L_i^0 \rangle$ acquires a non-vanishing vacuum expectation value (vev) after the electroweak breaking in a general basis. In our work, we adopt a basis such that the $\langle L_i^0 \rangle$ vanishes by a proper redefinition of $L_i$ and $H_d$. In this basis, the trilinear $R$-parity violating interactions are generically induced. Furthermore, the right-handed sneutrinos get mixed with the left-handed sneutrinos as well as up- and down-type Higgs, which would make analysis on the right-handed sneutrino decay complicated. The purpose of this paper is not to explore all possible parameter space, but to illustrate our basic idea that the dark matter with a lepton number can account for the sharp rise in the positron fraction while the antiproton flux can be suppressed. Thus we focus on a case that the dark matter is comprised of $N_3$, which decays into the SM particles through the $R$-parity violating bilinear term with $\mu_1 \neq 0$. Then, the relevant mixings of $N_3$ with the left-handed sneutrinos and the Higgs bosons are suppressed by the small mixing angle $\theta_{13}$. In particular, those mixings are absent in the
limit of \( \theta_{13} = 0 \), which therefore greatly simplifies our analysis. The adopted values of the mass differences and the mixing angles are shown in Table I. We will discuss later how our result is modified for other choices of the parameters.

In our set-up, the main decay channels of the right-handed sneutrino, \( \tilde{\nu}_R^3 \) (the scalar component of \( \tilde{N}_3 \)), are neutrinos and charged leptons through the neutrino-neutralino and charged-lepton-chargino mixings. The corresponding Feynman diagrams are shown in Fig. I. We can see from Eq. (3) that the interactions between the right-handed sneutrino \( \tilde{\nu}_R \) with the higgsino and the SM leptons are given by

\[
\mathcal{L} \supset -y_{ij} \tilde{\nu}_R^c \ell^c_i \ell^c_j H_u^0 + y_{ij} \tilde{\nu}_R^c \ell^c_j H^0_u,
\]

where \( \tilde{\nu}_R^c \) denotes the scalar component of \( \tilde{N}_i \), and \( \ell_1 = e, \ell_2 = \mu \) and \( \ell_3 = \tau \). First let us consider the neutrino-neutralino mixing. The mass matrix of the neutral fermions \( M_N \) is given by

\[
M_N = \begin{pmatrix} M^{(N)}_{MSSM} & M^{(N)}_R \\ M^{(N)\text{T}}_R & 0 \end{pmatrix},
\]

where we have defined \( \Psi^0 \equiv (\tilde{\nu}, \tilde{W}^3, \tilde{H}_d^0, \tilde{H}_u^0, \nu_{Li})^T \), \( M^{(N)}_{MSSM} \) is the usual neutralino mass matrix in MSSM, and \( M^{(N)}_R \) is the mass matrix between the neutralinos and the left-handed neutrinos arising from the \( R \)-parity violating interactions. The explicit expressions of \( M^{(N)}_{MSSM} \) and \( M^{(N)}_R \) are

\[
M^{(N)}_{MSSM} = \begin{pmatrix} M_1 & 0 & -m_Z s_W c_\beta & m_Z s_W s_\beta \\ 0 & M_2 & m_Z c_W c_\beta & -m_Z c_W s_\beta \\ -m_Z s_W c_\beta & m_Z c_W c_\beta & 0 & -\mu \\ m_Z s_W s_\beta & -m_Z c_W s_\beta & -\mu & 0 \end{pmatrix}, \quad M^{(N)}_R = \begin{pmatrix} 0_{3 \times 3} \\ -\mu_1 & -\mu_2 & -\mu_3 \end{pmatrix},
\]

where \( c_W \equiv \cos \theta_W \) and \( s_W \equiv \sin \theta_W \) with \( \theta_W \) being the weak mixing angle; \( c_\beta \equiv \cos \beta \) and \( s_\beta \equiv \sin \beta \) with \( \tan \beta = v_u/v_d \) where \( v_u(v_d) \) is the vev of up-type (down-type) Higgs field. Here we keep \( \mu_2 \) and \( \mu_3 \) in the mass matrix, although we will set them to be zero in the following analysis. The mass matrix \( M_N \) could be diagonalized by an unitary matrix \( V \), and the mixing between the up-type higgsino \( \tilde{H}_u^0 \) and the left-handed neutrinos \( \nu_{Li} \) is given by \( V_{4,i+4} \). In a similar way, the charged leptons mix with the charginos via the \( R \)-parity
FIG. 1: Feynman diagrams of right-handed sneutrino decay via the mixing between the up-type higgsino and the SM leptons.

violating interactions. The mass matrix $M_C$ of the charginos and the charged leptons is given by $\mathcal{L} \supset -\Psi^{-T} M_C \Psi^+$ with \[26\]

$$M_C = \begin{pmatrix} M_{MSSM}^{(C)} & 0_{2 \times 3} \\ M_{R}^{(C)} y_i^e \delta_{ij} v_d / \sqrt{2} \end{pmatrix},$$

where we have defined $\Psi^{-} \equiv (\tilde{W}^{-}, \tilde{H}_{d}^{-}, \ell_i)^T$ and $\Psi^{+} \equiv (\tilde{W}^{+}, \tilde{H}_{u}^{+}, e_i^c)^T$. $M^{(C)}_{MSSM}$ is the usual chargino mass matrix in MSSM, $M_{R}^{(C)}$ is the mass matrix induced by $R$-parity violating interactions and $y^e$ is the charged lepton Yukawa coupling. The explicit expressions are

$$M_{MSSM}^{(C)} = \begin{pmatrix} M_2 & \sqrt{2} m_Z c_W s_\beta \\ \sqrt{2} m_Z c_W c_\beta & \mu \end{pmatrix}, \quad M_{R}^{(C)} = \begin{pmatrix} 0 & -\mu_1 \\ -\mu_2 & 0 \\ -\mu_3 & 0 \end{pmatrix}$$\(12\)

The mass matrix $M_C$ can be diagonalized by two rotational matrices $O_-$ and $O_+$ which relate the gauge eigenstates of $\Psi^{-}$ and $\Psi^{+}$ with their mass eigenstates, respectively, i.e. $O^T M_C O_\pm = \text{diag}(m_{\tilde{\chi}^{-}}, m_{\tilde{\chi}^{0}}, m_e, m_\mu, m_\tau)$. The mixings between $\tilde{H}_u^+$ and the right-handed charged leptons $e_i^{c+}$ are given by the elements $[O_+]_{2,2+i}$. Combining the interaction (7) and the mixings between the up-type higgsino and the SM leptons, the $\tilde{\nu}_R$ will decay into the SM leptons as shown in Fig. 1.

Finally, we summarize the relevant decay processes for our study below;

$$\tilde{\nu}_{Ri} \rightarrow \nu_j \bar{\nu}_k \propto |y^\nu_{ij} [V]_{4,k+4}|^2,$$

$$\tilde{\nu}_{Ri} \rightarrow \ell_j^- e_k^{c+} \propto |y^\nu_{ij} [O^+]_{2,2+k}|^2,$$

where $V$ and $O_\pm$ are the rotation matrices in the neutralino-neutrino and the chargino-lepton respectively. In Table II we show the branching ratios of a decaying $\tilde{\nu}_{R3}$, in which we assume non-vanishing $\mu_1$ with $\mu_{2,3} = 0$ for simplicity. Notice that positrons are always
produced in the charged lepton decay modes since only $\mu_1$ is non-vanishing. Also, because of the smallness of the positron mass, the decay branching ratios of positron are highly suppressed down to a few percentages, compared to that of the neutrino production.

III. COSMIC-RAY FLUXES

Let us here mention that both the right-handed sneutrino and its antiparticle can be dark matter and contribute to the cosmic-ray signals. For simplicity, we assume that both of them have been produced with an equal amount in the early universe, therefore, we have

$$\Omega_{DM} \equiv \frac{\rho_{\tilde{\nu}_R} + \rho_{\tilde{\nu}_c}}{\rho_c} = 2 \frac{\rho_{\tilde{\nu}_R}}{\rho_c} = 2 \Omega_{\tilde{\nu}_R},$$

where $\Omega_{DM}$ and $\Omega_{\tilde{\nu}_R}$ are density parameters of the dark matter and the right-handed sneutrino, respectively; $\rho_c$ is the critical density and $\rho_{\tilde{\nu}_R(\tilde{\nu}_c)}$ is the energy density of right-handed (anti)sneutrino. In the rest of this paper, contributions of both right-handed sneutrino and its antiparticle are included in the cosmic-ray fluxes. The method for calculations of the gamma-ray flux and positron fraction are the same as that in our previous study [9], therefore, we only show the equations which are needed in the calculations. We refer readers who are interested in the derivations to Ref. [7] and references therein.

A. Positron fraction

After being produced from the decay of the right-handed sneutrino, positrons will propagate in the magnetic field of the Milky Way. The positron flux is given by

$$\Phi_{e^+}^{prim}(E) = \frac{c}{4\pi m_{\tilde{\nu}_R} \tau_{\tilde{\nu}_R}} \int_0^{m_{\tilde{\nu}_R}/2} dE' G(E, E') \frac{dN_{e^+}}{dE'} ,$$

where $E$ is in units of GeV, and $m_{\tilde{\nu}_R}$ and $\tau_{\tilde{\nu}_R}$ are the mass and lifetime of the right-handed sneutrino dark matter. Here $dN_{e^+}/dE$ is the energy spectrum of positron from the dark matter decay. We have used PYTHIA [27] to calculate the energy spectra, and the numerical

| channel | $e^+\mu^-$ | $e^+\tau^-$ | $\nu\bar{\nu}$ |
|---------|-------------|-------------|-----------------|
| branching ratios (%) | 1.5 | 1.5 | 97 |

TABLE II: Branching ratios of right-handed sneutrino $\tilde{\nu}_{R3}$. 
results are shown in Fig. 2. The Green function $G(E, E')$ in Eq. (15) is approximately given by 

$$G(E, E') \simeq \frac{10^{16}}{E^2} e^{a+b(E^{\delta -1} - E'^{\delta -1})} \theta(E' - E) \quad [\text{sec/cm}^3],$$

(16)

where $\delta$ is related to the properties of the interstellar medium and can be determined mainly from the Boron to Carbon ratio ($B/C$) [28]. In the numerical study, we adopt the following parameters, $\delta = 0.55$, $a = -0.9716$ and $b = -10.012$ [7], that are consistent with the B/C value and produce the minimum flux of positrons.

In addition to the positron flux from dark matter decay, there exist secondary positrons produced from interactions between the primary cosmic rays and nuclei in the interstellar medium. The positron flux is considered to suffer from the solar modulation, especially for the energy below 10 GeV. If the solar modulation effect is independent of the charge-sign, one can cancel the effect by taking the positron fraction,

$$\frac{\Phi_{e^+}}{\Phi_{e^+} + \Phi_{e^-}},$$

(17)

which is indeed measured in many experiments. To estimate the positron fraction, it is necessary to include the electron flux. We use the approximations of the $e^-$ and $e^+$ background...
FIG. 3: The fraction of positron flux from right-handed sneutrino dark matter decay, shown together with experimental data [13, 21, 22, 46].

The typical feature is the sharp turnup at $E \approx 10$ GeV and a drop-off at $E = m_{\tilde{\nu}_R}/2$ due to the contribution of the direct production of $e^+$ from the decay of dark matter [9].

B. Gamma-ray flux

Generally, the main contribution to the gamma-ray spectrum arises from the pion $\pi^0$ generated in the QCD hadronization process of $q\bar{q}$ and in the decay of $\tau$. Since no quark
pair is produced in our set-up, the only source of $\pi^0$ is from the decay of $\tau$. To estimate the spectrum, we use the PYTHIA \cite{27} Monte Carlo program with the branching ratios shown in Table \textbf{II} and the energy spectra $dN_\gamma/dE$ are presented in Fig. \textbf{2}. It is worth noting that there is no line emission of the gamma rays from the decay of $\tilde{\nu}_R$, which is present in the case of the gravitino dark matter \cite{7, 8}.

There are galactic and extragalactic contributions from the decay of $\tilde{\nu}_R$ to the observed gamma-ray flux. The flux of the gamma-ray from the extragalactic origin is estimated as

\begin{equation}
\frac{E^2 dJ_\gamma}{dE} = \frac{E^2 c \Omega_{DM} \rho_c}{4 \pi m_{\tilde{\nu}_R} \tau_{\tilde{\nu}_R} H_0 \Omega_M^{1/2}} \int_1^{y_{eq}} dy \frac{dN_\gamma}{d(yE)} \frac{y^{-3/2}}{\sqrt{1 + \Omega_M y^{-1} \Omega_{\Lambda} y^{-3}}}, \tag{20}
\end{equation}

where $c$ is the speed of light; $\Omega_M$ and $\Omega_{\Lambda}$ are the density parameters of matter (including both baryons and dark matter) and the cosmological constant, respectively; $H_0$ is the Hubble parameter at the present time; $y \equiv 1 + z$, where $z$ is the redshift, and $y_{eq}$ denotes a value of $y$ at the matter-radiation equality. For the numerical results, we use \cite{1}

\begin{align*}
\Omega_{DM}h^2 &= 0.1099, \quad \Omega_M h^2 = 0.1326, \quad \Omega_{\Lambda} = 0.742, \quad \rho_c = 1.0537 \times 10^{-5} \text{GeV/cm}^3. \tag{21}
\end{align*}

On the other hand, the gamma-ray flux from the decay of dark matter in the Milky Way halo is

\begin{equation}
\frac{E^2 dJ_\gamma}{dE} = \frac{E^2}{4 \pi m_{\tilde{\nu}_R} \tau_{\tilde{\nu}_R} H_0} \int_{\rho_{\text{los}}} \rho_{\text{halo}}(\vec{r}) d\vec{r}, \tag{22}
\end{equation}

where $\rho_{\text{halo}}$ is the density profile of dark matter in the Milky Way, $\int_{\text{los}} \rho_{\text{halo}}(\vec{r}) d\vec{r}$ is the average of the integration along the line of sight (los). In our calculation, we adopt the Navarro-Frenk-White (NFW) halo profile \cite{32}

\begin{equation}
\rho(r) = \frac{\rho_0}{(r/r_c)(1 + r/r_c)^2}, \tag{23}
\end{equation}

where $r$ is the distance from the center of Milky Way, $r_c = 20$ kpc, and $\rho_0$ is set in such a way that the dark matter density in the solar system satisfies $\rho(r_\odot) = 0.30$ GeV/cm$^3$ \cite{33} with $r_\odot = 8.5$ kpc being the distance from the Sun to the Galactic Center. For the background, we use a power-law form adopted in Ref. \cite{8}

\begin{equation}
\frac{E^2 dJ_\gamma}{dE} \approx 5.18 \times 10^{-7} E^{-0.499} \text{ GeVcm}^{-2}\text{sr}^{-1}\text{sec}^{-1}, \tag{24}
\end{equation}

where $E$ is in units of GeV. The predicted gamma-ray flux from the decaying $\tilde{\nu}_R$ are shown in Fig. \textbf{4} together with experimental data. Since no hadronic decay mode exists and the
amount of $\tau$ is insignificant, the predicted gamma-ray flux is below the EGRET data for the parameters we have chosen. However, we do not regard this tension as a serious conflict, and we expect that the upcoming data from FGST should be able to further probe such a prediction from the right-handed sneutrino dark matter decay. Also, we should emphasize that the relatively suppressed gamma-ray flux shown in Fig. 4 is due to the simplification we have adopted. The gamma-ray flux may fit the EGRET data if we allow $\mu_{2,3}$ to be non-zero and/or the other generations of the right-handed sneutrino to contribute to the dark matter abundance. This is because the decay branching ratios of $\tau$ will increase and $q\bar{q}$ decay channel will open due to the mixing between the right-handed sneutrinos and the Higgs bosons, then obviously, more $\pi^0$ will be generated.

C. Antiproton flux

As we have mentioned before, the predicted antiprotons from decay of gravitino dark matter tend to have a tension with the observational data. Of course, this apparent tension may not be serious at all, considering the large uncertainties on both experimental data and diffusion models. However, we should better have a scenario in which the antipro-
ton flux could be suppressed. In our set-up, the right-handed sneutrino does not decay into quarks (at least the decay rate is suppressed), and therefore the antiproton flux is negligible. This feature will be lost once we allow $\theta_{13}$ and/or $\mu_{2,3}$ to deviate from zero. Or if $\tilde{\nu}_{R1,2}$ also contribute to the dark matter abundance, their decay generically produces the antiproton, since they can mix with the Higgs bosons as well as the left-handed sneutrinos. The antiproton flux in generic parameter space will depend on those parameters as well as the neutrino mass spectrum, and it is beyond the scope of this paper to survey all the possibilities. However, it is important to keep in mind that there is a set of parameters where the antiproton production becomes negligible.

D. Neutrino flux

The main decay channel of the $\tilde{\nu}_{R3}$ is $\nu\bar{\nu}$ as can be seen from Table II. It has been recently studied in detail whether the neutrinos produced by the gravitino decay can be detected by the current and future observations [34]. In our scenario, the neutrino flux is enhanced compared to the positron and the gamma-ray ones, and therefore the neutrino signal from the $\tilde{\nu}_R$ decay can be larger by an order of magnitude compared to the gravitino of the same mass. Thus, the neutrino signal, especially $\nu_\tau$, may reach the sensitivity of future experiments, e.g. Hyper-Kamiokande. However, we need to mention that the enhancement in the neutrino production will be lost for another choice of the parameters.

IV. COSMOLOGICAL PRODUCTION OF $\tilde{\nu}_R$

We would like to briefly discuss here how the right-handed sneutrino is produced in the early universe. In order to account for the observed dark matter abundance, the cosmological abundance of $\tilde{\nu}_R$ must satisfy [1]

$$\Omega_{\tilde{\nu}_R} h^2 + \Omega_{\tilde{\nu}_c} h^2 = 0.1099 \pm 0.0062,$$

or equivalently,

$$\frac{n_{\tilde{\nu}_R} + n_{\tilde{\nu}_c}}{s} \simeq 4.2 \times 10^{-12} \left( \frac{m_{\tilde{\nu}_R}}{100 \text{GeV}} \right)^{-1},$$

where $n_{\tilde{\nu}_R}$ is the number density of $\tilde{\nu}_R$, and $s$ the entropy density of the universe. If not only the lightest but also heavier right-handed sneutrinos are stable in a cosmological time, we need to sum over those right-handed sneutrinos in Eq. [26].
In the case of the neutralino dark matter, it is usually assumed that the neutralinos are thermally produced at a temperature above the freeze-out temperature. However, since the neutrino Yukawa coupling $y^\nu$ is very small ($\sim 10^{-13}$) in our scenario, the $\tilde{\nu}_R$ is never in equilibrium unless it has other unknown strong interactions with the radiation in the early universe. Therefore we need to consider non-thermal production of $\tilde{\nu}_R$. There are several possibilities [18], and we consider one by one as follows.

First of all, $\tilde{\nu}_R$ is produced by the decay of the lightest SUSY particle in the MSSM sector (MSSM-LSP) through the neutrino Yukawa coupling(s). In the presence of $R$-parity violating interactions, the MSSM-LSP also decays into the SM particles without producing $\tilde{\nu}_R$. If the $R$-parity violating coupling is much larger than the neutrino Yukawa coupling, the production of $\tilde{\nu}_R$ will get suppressed. This is the case in our set-up, since a relatively large $R$-parity violation is needed to compensate the suppression of the decay rate into the charged leptons. However, for another choice of the model parameters, the magnitude of the $R$-parity violation tends to be as small as the neutrino Yukawa coupling, and the MSSM-LSP decay into $\tilde{\nu}_R$ via the $R$-parity conserving interactions may occur at a non-negligible rate. This would be indeed the case if the neutrino masses are degenerate, since the neutrino Yukawa coupling will become larger than that in the case of normal hierarchy, and the needed coupling strength of the $R$-parity violation to account for the positron excess becomes correspondingly smaller. Then some amount of $\tilde{\nu}_R$ will be produced by the decay of the MSSM-LSP. In the $R$-parity conserving case, the production from the MSSM-LSP was studied in Ref. [18]. We can estimate the abundance in a similar way, taking account of the suppression due to the presence of the decay via $R$-parity violating interactions. That is to say, contribution to the density parameter of $\tilde{\nu}_R$ from the decay of the MSSM-LSP, $\chi$, is given by

$$\Delta \Omega_{\tilde{\nu}_R} = B_\nu \left( \frac{m_{\tilde{\nu}_R}}{m_\chi} \right) \Omega_\chi,$$

where $B_\nu$ denotes the branching ratio of the $\chi$ decay into $\tilde{\nu}_R$, and $m_\chi$ is the mass of $\chi$. Note that, due to the presence of the $R$-parity violation, the $\chi$ decays faster than the case without it, which makes it easier to satisfy the big bang nucleosynthesis (BBN) bound on the MSSM-LSP decay. However, it also means that larger abundance of the MSSM-LSP is needed to produce a right amount of $\tilde{\nu}_R$. In the normal hierarchy case, this would push up the soft masses of the SUSY particles, and the naturalness may be called into question. On the other hand, in the case of degenerate neutrino masses, the effect of the $R$-parity

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violation is relatively smaller, and we expect that a right amount of $\tilde{\nu}_R$ can be generated by the MSSM-LSP decay.

Another interesting possibility is that the gravitino decay produces $\tilde{\nu}_R$. The gravitinos are generated by particle scattering in thermal plasma [35, 36], and they are also generically produced by the inflaton decay [37, 38, 39, 40, 41]. For the gravitino mass is of $\mathcal{O}(100)$ GeV, it typically decays into the SM particles and their superpartners during BBN. The energetic decay products can significantly change abundances of light elements, thus the gravitino abundance is tightly constrained by BBN. If the gravitino is the next-to-lightest SUSY particle (NLSP), however, it will mainly decay into $\tilde{\nu}_R$ and a neutrino, and therefore, the BBN constraint will be greatly relaxed. To account for the abundance Eq. (26), the reheating temperature should be as high as about $10^{10}$ GeV if the gravitino is mainly thermally produced. Note that the production of $\tilde{\nu}_R$ from the gravitino decay can be concomitant with the production from the MSSM-LSP decay as well as the following mechanisms, so the reheating temperature can be lower.

The inflaton decay may also directly produce $\tilde{\nu}_R$. Indeed, if the inflaton acquires a non-vanishing vev at the potential minimum, it couples to all the matter fields as well as the gauge fields in supergravity [39, 40, 41]. However, the smallness of the neutrino Yukawa coupling makes the branching ratio of the decay into $\tilde{\nu}_R$ extremely small. It is likely that we need to assume relatively strong interactions between $\tilde{\nu}_R$ and the inflaton sector (or other fields whose energy dominates the energy density of the universe after inflation) for the enhancement of contribution from inflaton decay.

The last possibility is the scalar condensation of $\tilde{\nu}_R$. If the position of $\tilde{\nu}_R$ is deviated from the origin during inflation $^2$, $\tilde{\nu}_R$ will start oscillating after inflation when the Hubble parameter becomes comparable to its mass. The abundance is estimated to be

$$\frac{n_{\tilde{\nu}_R}}{s} \simeq \frac{T_R}{4m_{\tilde{\nu}_R}} \left( \frac{\tilde{\nu}_R^{osc}}{M_P} \right)^2 \sim 2 \times 10^{-13} \left( \frac{m_{\tilde{\nu}_R}}{100 \text{ GeV}} \right)^{-1} \left( \frac{T_R}{10^6 \text{ GeV}} \right) \left( \frac{\tilde{\nu}_R^{osc}}{10^{10} \text{ GeV}} \right)^2,$$

where $T_R$ denotes the reheating temperature of the universe and $\tilde{\nu}_R^{osc}$ is the amplitude of the oscillations. We have assumed that the sneutrino starts to oscillate before the reheating is completed. For an appropriate choice of the amplitude, it can account for the right

$^2$ If $\tilde{\nu}_R$ is light during inflation, such a displacement is generically expected to arise from the quantum fluctuations of the $\tilde{\nu}_R$. The fluctuations will become isocurvature fluctuations in the CDM. It is also possible to generate large amount of non-Gaussianity in the CDM isocurvature perturbations [42].
abundance of the observed dark matter. Note that the $\Omega_{R}$ is determined solely by the reheating temperature and the initial position of $\tilde{\nu}_R$ in this case, and is independent of the sneutrino mass.

V. DISCUSSION AND CONCLUSIONS

We have assumed that the $R$-parity is dominantly violated by the bilinear term. If it is violated mainly by the trilinear terms, $LL\tilde{E}$, $Q\tilde{D}L$, or $\tilde{U}\tilde{D}\tilde{D}$, the predicted cosmic-ray spectra will be different. If the $R$-parity is broken by the $LL\tilde{E}$ operator, the $\tilde{\nu}_R$ will decay into charged leptons as well as neutrinos through the left-right mixing, which will result in a sharp peak in the predicted positron fraction at the solar system. The continuum gamma rays are also produced, while virtually no antiprotons are produced. This choice therefore provides an ideal way to avoid the observational constraint on the antiproton flux. In the case of $Q\tilde{D}L$ or $\tilde{U}\tilde{D}\tilde{D}$, the hadronic branching ratios are larger than the case of $LL\tilde{E}$, and therefore we expect a relatively large contribution to the antiproton flux.

With the tiny $R$-parity violations in our scenario, the NLSP is generally long-lived in collider experiments and decays outside the detector \cite{43, 44}. If NLSP is the neutralino, it is observed as a missing energy, and the collider phenomenology is the same as the $R$-parity conserved case. If stau is the NLSP, we expect to observe its track inside the collider, which is similar to the case of the decaying gravitino LSP \cite{8}. However, there is one possibility that the $\tilde{\nu}_R$ is only a fraction of the total dark matter, and the rest is explained by some other stable particles, such as a QCD axion. If the fraction of $\tilde{\nu}_R$ in dark matter $r$ is very small, the $R$-parity violating coupling should be enhanced as $\propto r^{-1/2}$, for a fixed positron flux. Since the lifetime must be longer than the present age of the universe, $r$ can be as small as $O(10^{-10})$, and the corresponding $R$-parity violation can be enhanced by a factor of $O(10^5)$. With such a large enhancement in the $R$-parity violation, we may observe some signatures of decaying NLSP at colliders. Note also that, even if the fraction is much smaller than unity, the features of the cosmic-ray spectra from decaying right-handed sneutrino will not change.

We have focused on the case that all the right-handed sneutrinos have almost the same mass, and considered only one of them is the source of cosmic-ray. If the three right-handed sneutrinos exist with an equal amount in our universe, they will all contribute to the cosmic-
ray. The contribution from one of the three right-handed sneutrinos may dominate over those from the other two, depending on the values of $\mu_i$ and the neutrino mass spectrum.

Even there exist mass differences among the three right-handed sneutrinos, unless the decay channel for the heavier one to a on-shell higgsino is kinematically allowed, the decay patterns of all the right-handed sneutrinos are similar to each other. Namely, the heavier $\tilde{\nu}_R$ also decays via $R$-parity violating interactions, since the $R$-parity conserving decay channel is suppressed as explained in the following paragraph.

We should also emphasize here that the excess in the cosmic-ray positron can be also explained without introducing the $R$-parity violation interactions. If the right-handed sneutrinos are the three lightest supersymmetric particles, and if there is slight hierarchy in the masses, the heavier one will decay into a lighter one accompanied by two SM leptons through an off-shell higgsino. It is obvious that the two SM leptons will be the source for the cosmic positron and gamma-ray excesses. Also, the antiproton will not be produced. However, the lifetime tends to be much longer than the needed one to account for the positron excess, since the decay amplitude is suppressed by the neutrino Yukawa coupling squared. One solution is to add a small Majorana mass term for the right-handed neutrinos [19]. Then the see-saw mechanism will occur at very low energy, and the neutrino Yukawa coupling becomes larger, which can make the lifetime of the heavier one short enough to explain the positron excess. Another way to make the lifetime shorter is to tune the higgsino mass close to the heavier right-handed sneutrino mass.

It is even possible to explain the positron excesses in a non-supersymmetric theory. One candidate is the sterile neutrino dark matter, whose decay produces charged leptons and neutrinos. The decay processes of a sterile neutrino lighter than 100 MeV were studied in a different context (see e.g. [45]). Since the sterile neutrino must be heavier than 100 GeV in our case, it will decay also into $W$ and $Z$ bosons, and also into $\nu\gamma$ with a suppressed rate. Thus, the energy spectra of the cosmic-ray particles from the sterile neutrino decay will look similar to those in the case of the decaying gravitino. The mixing angle between the sterile and active neutrinos must be extremely suppressed in order to realize the long lifetime of the sterile neutrino, which will make it difficult to produce the sterile neutrino through the mixing. It will be necessary to introduce a non-thermal production mechanism of the sterile neutrino.

So far, only the normal hierarchy case in neutrino mass spectrum is considered, but we
expect that an essential feature of our results will persist in the other neutrino mass spectra. Rather, it may be possible to extract some information on the neutrino mass spectrum from the positron and gamma-ray fluxes by the PAMELA and FGST in operation. We leave those issues for future work.

We propose a leptonic dark matter candidate whose decay may explain the excesses of observed cosmic rays. If the dark matter has a lepton number, it will not be surprising that we see some excesses only in the positron and the gamma-ray fluxes while virtually no primary antiprotons are observed. To illustrate the idea, in this paper we have studied a scenario that the neutrinos are Dirac type, and the right-handed sneutrino accounts for the observed dark matter and decays into the SM particles through the $R$-parity violating interactions. The charged leptons as well as neutrinos are directly produced during the decay, leading to a sharp rise in the positron fraction. The decay products will also generate continuum gamma rays, which may account for the EGRET excess. Interestingly, with an appropriate set of the parameters, the antiproton production can be suppressed in our scenario, which is consistent with an observation that most of the observed antiprotons are secondaries. Furthermore, we expect that the energy spectra for those cosmic-ray particles may provide a new probe into the neutrino mass spectrum. Those features on the predicted cosmic-ray spectra in our study will be checked by the PAMELA and FGST satellites.

**Note added:** After our paper was posted, the PAMELA group reported a steep rise in the positron fraction [46], which can be explained by our model with a 300 GeV right-handed sneutrino dark matter, as shown in Fig. 3.

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