Decaying LSP in SO(10) GUT and 
PAMELA’s Cosmic Positron

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Abstract

We suppose that the lightest supersymmetric particle (LSP) in the minimal supersymmetric
standard model (MSSM) is the dark matter. The bino-like LSP can decay through the SO(10)
gauge interactions, if one right-handed (RH) neutrino ($\nu_1^c$) is lighter than the LSP and its superpartner ($\tilde{\nu}_1^c$) develops a vacuum expectation value (VEV), raising extremely small R-parity violation naturally. The leptonic decay modes can be dominant, if the VEV scale of $16_H$ is a few orders of magnitude lower than the VEV of $45_H$ ($\approx 10^{16}$ GeV), and if a slepton ($\tilde{e}_1^c$) is relatively lighter than squarks. The desired decay rate of the LSP, $\Gamma_\chi \sim 10^{-26}$ sec.$^{-1}$ to explain PAMELA data can be naturally achieved, because the gaugino mediating the LSP decay is superheavy. From PAMELA data, the $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ breaking scale (or the $16_H$ VEV scale) can be determined. A global symmetry is necessary to suppress the Yukawa couplings between one RH (s)neutrino and the MSSM fields. Even if one RH neutrino is quite light, the seesaw mechanism providing the extremely light three physical neutrinos and their oscillations is still at work.

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I. INTRODUCTION

For the last three decades, many remarkable progresses in particle physics and cosmology have been made thanks to the cooperative and intimate relation between the two fields. In particular, the application of particle physics theory into dark matter (DM) models in cosmology was very successful. Because of the correct order of magnitude of the cross section, thermally produced weakly interacting massive particles (WIMPs) have been long believed to be DM candidates [1]. So far the lightest supersymmetric particle (LSP), which is a well-motivated particle originated from the promising particle physics model, i.e. the minimal supersymmetric standard model (MSSM), has attracted much attentions as an excellent example of WIMP.

Recently, PAMELA [2], ATIC [3], H.E.S.S. [4], and the Fermi-LAT collaborations [5] reported the very challenging observations of positron excesses in cosmic ray above 30 GeV up to the TeV scale. In particular, PAMELA observed a positron fraction \(e^+/(e^+ + e^-)\) exceeding the theoretical expectation [6] above 30 GeV up to 100 GeV. However, the antiproton/proton flux ratio was quite consistent with the theoretical calculation. The ATIC, H.E.S.S., and Fermi-LAT’s observations exhibit excesses of \((e^+ + e^-)\) flux in cosmic ray from 100 GeV to 1 TeV.\(^1\) They would result from the positron flux that keeps rising up to 1 TeV.

Apparently the above observational results are very hard to be interpreted in view of the conventional MSSM cold dark matter scenario: explaining the excess positrons with annihilations of Majorana fermions such as the LSP needs a too huge boost factor. Moreover, ATIC, H.E.S.S., and Fermi-LAT’s observations seem to require a TeV scale DM, if they are caused indeed by DM annihilation or decay. Introduction of a TeV scale LSP, however, would spoil the motivation of introducing supersymmetry (SUSY) to resolve the gauge hierarchy problem in particle physics. In addition, TeV scale DM seems to be disfavored by the gamma ray data [4], if the excess positron flux is due to DM annihilations [7]. On the other hand, the DM decay scenario is relatively free from the gamma ray constraint [8].

In the DM decay scenario, however, there are some serious hurdles to overcome: one is to naturally obtain the extremely small decay rate of the DM \((\Gamma_{DM} \sim 10^{-26} \text{ sec}^{-1})\), and the other is to naturally explain the relic density of the DM in the Universe. The first hurdle

\(^1\) H.E.S.S. measured Cherenkov radiations by cosmic electrons and positrons above 600 GeV energy scale.
could be somehow resolved by introducing an extra symmetry, an extra DM component with a TeV scale mass, and grand unified theory (GUT) scale superheavy particles, which mediate DM decay into the SM charged leptons (and the LSP) [9]. The fact that the GUT scale particles are involved in the DM decay might be an important hint supporting GUT [10, 11]. However, since the interaction between the new DM and the SM charged lepton are made extremely weak by introducing superheavy particles mediating the DM decay, non-thermal production of the DM with a carefully tuned reheating temperature should be necessarily assumed. One way to avoid it is to consider SUSY models with two DM components [9, 10]. In these models, the decay of the small amount of the meta-stable heavier DM component \((X)\), which is assumed to be non-thermally produced, accounts for the cosmic positron excess, and the thermally produced lighter DM component LSP \((\chi)\), which is absolutely stable and regarded as the dominant DM \([\mathcal{O}(10^{-10}) < n_X/n_\chi]\), explains the relic density of the Universe.\(^2\)

In this paper, we suppose that the conventional bino-like LSP is the main component of the DM. Since the “bino” is a WIMP, thermally produced bins could explain well the relic density of the Universe. The bino-like LSP with a mass of about 300 – 400 GeV could also explain PAMELA data, if it decays to \(e^\pm\) and a neutral fermion with an extremely small decay rate of order \(10^{-26}\) sec.\(^{-1}\) [13]. The \((e^+ + e^-)\) excess observed by Fermi-LAT could be explained by astrophysical sources such as nearby pulsars [14] (and/or with the sub-dominant extra TeV scale DM component [9]).\(^3\) In fact, pulsars can explain both the PAMELA and Fermi-LAT’s data in a suitable parameter range [15]. However, this does not imply that DM in addition to pulsars can not be the source of the galactic positrons [14]. In fact, we don’t know yet a complete pulsar model, in which all the free parameters would be fixed by the fundamental physical constants.

To achieve the needed extremely small decay rate of the bino-like LSP \(\chi\), we need extremely small R-parity violation. We will assume that the R-parity is broken by a non-zero vacuum expectation value (VEV) of a right-handed (RH) sneutrino \((\langle \tilde{\nu}_1 \rangle \neq 0)\). Since it

\(^{2}\) The low energy field spectrum in the models of Ref. [9] is the same as that of the MSSM except for the neutral singlet extra DM component. Moreover, the models in [9] can be embedded in the flipped SU(5) GUT and string models [11, 12].

\(^{3}\) Alternatively, one could assume a bino mass of 3.5 TeV in order to account for both PAMELA and Fermi-LAT with LSP decay [13]. In this case, however, the soft SUSY breaking scale should be higher than 3.5 TeV.
doesn’t carry any standard model (SM) quantum number, it does not interact with the MSSM fields at all, if its Yukawa interactions with them are forbidden by a symmetry and gravity interaction is ignored. We will explore the possibility that the extremely small DM decay rate results from the gauge interaction by exchange of the superheavy gauge bosons and gauginos present in the SO(10) SUSY GUT. We will not introduce a new DM component, and will attempt to explain the PAMELA’s observation within the framework of the already existing particle physics model.

II. SO(10) GUT

One of the appealing GUTs is the SO(10) GUT [16]. It unifies all the three SM gauge forces within the SO(10) gauge interaction. One of the nice features of SO(10) is that it predicts the existence of the RH neutrinos [or the SU(2)$_L$ singlet neutrinos], since a RH neutrino is contained in a single spinorial representation $16$ of SO(10), together with one family of the SM fermions. The RH neutrinos provide a very nice explanation of the observed neutrino oscillations through the seesaw mechanism [17] and also of the baryon asymmetry in the Universe through leptogenesis [18].

A. Superheavy fields in SO(10)

SO(10) GUT models contain many superheavy particles. They might be utilized to get the required DM decay rate of $10^{-26}$ sec.$^{-1}$ Most of all, the gauge bosons and gauginos corresponding to the coset SO(10)/SM have masses around the GUT scale. In this paper, we are particularly interested in them as the mediators of DM decay.

The superfields in the Higgs sector needed for breaking SO(10) to the SM are also superheavy. Particularly, an adjoint Higgs $45_H$ (or $210_H$) and spinorial Higgs $16_H$ and $\overline{16}_H$ (or $126_H$ and $\overline{126}_H$) can be employed to achieve the SM gauge group from SO(10). The vector representation $10_h$, which includes the two MSSM Higgs doublets, contains also the superheavy Higgs triplets $D_H$ and $D'_H$. Their masses can be obtained through a proper doublet/triplet splitting mechanism. One way is to introduce the coupling $10_h \cdot 45_H \cdot 10_h$, assuming the scalar component of $45_H$ develops a VEV along the SU(3)$_c \times$SU(2)$_L \times$SU(2)$_R \times$U(1)$_{B-L}$
(≡ LR) direction \[19\]. In this paper, we will thus identify the triplet Higgs mass scale with the VEV of $45_H$.

How many and what kind of Higgs fields are needed to get the SM gauge group are quite model-dependent. Their masses would be close to the GUT scale, but they are not exactly the same as each other. Even in one Higgs multiplet, its component fields might have various mass spectra after symmetry breaking. Except for $10_h(45_H)10_h$, they interact with the MSSM fields only through non-renormalizable Yukawa couplings due to their GUT scale VEVs. Such couplings can be utilized to get the realistic SM fermion masses. One might think that SO(10)-breaking superheavy fields also contribute to the mediation of DM decay through such non-renormalizable couplings with the MSSM fields. However, the extra suppression factor $(1/M_P)^n$ ($n = 1, 2, 3, \cdots$) makes their contributions negligible compared to those of the superheavy gauge fields and gauginos via the renormalizable gauge interactions, which will be discussed later.

The SO(10)-breaking sector could include heavy fields, which do not develop GUT scale VEVs. They are introduced in order to decouple unwanted fields in the SO(10)-breaking Higgs sector, which are absent in the MSSM, from low energy physics in non-minimal SO(10) models. Since their couplings to the MSSM fields are not essential and their masses would be heavier than the mediators leading to DM decay, we can assume that all the interactions between such SO(10)-breaking sector fields and the MSSM fields are weak enough, if they are present.

Thus, as far as the DM decay is concerned, the gauge interactions through the superheavy gauge fields and gauginos can be dominant over Yukawa interactions. They would give more predictable results, regardless of what specific SO(10) models are adopted. We will focus on the DM decay predominantly through the superheavy gauge fields or gauginos.

B. $SU(5)$ vs. $SU(2)_R$ scale

In terms of the SM’s quantum numbers, the SO(10) generator (= $45_G$) is split into the SM gauge group’s generators plus \{(1, 1)_{-1}, (1, 1)_1\}, \{(1, 1)_0\}, \{(3, 2)_{-5/6}, (\bar{3}, 2)_{5/6}\}, and

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4 When SO(10) is broken by $16_H$, $\overline{16}_H$, and $45_H$, and the doublets/triplets in $10_h$ are split by the coupling $10_h(45_H)10_h$, the pseudo-goldstones included in the Higgs would not become easily superheavy. Then the Higgs sector needs to be extended by introducing more superfields and specific interactions \[20\].
\{(3,2)_{1/6}, (\bar{3},2)_{-1/6}; (3,1)_{2/3}, (\bar{3},1)_{-2/3}\}. We will simply write them as

\{E, E^c\}, \ N, \ \{Q', Q'^c\}, \ \{Q, Q^c; U, U^c\}, \ (1)

respectively. By the VEV of the adjoint Higgs \langle 45_H \rangle, the SO(10) gauge symmetry may break to LR. Through this process, the gauge boson and the gauginos carrying the quantum numbers of \{Q', Q'^c\} and \{Q, Q^c; U, U^c\} achieve heavy masses proportional to \langle 45_H \rangle. The \{E, E^c\} and a linear combination of the SM hypercharge generator and \ N (≡ N_R) composes the SU(2)_R generators. The other combination orthogonal to it corresponds to the U(1)_{B-L} generator (≡ N_{BL}). They don’t get masses from \langle 45_H \rangle.

On the other hand, the VEVs of the Higgs in the spinorial representations \langle 16_H \rangle, \langle \bar{16}_H \rangle breaks SO(10) down to SU(5). This process generates the heavy masses proportional to \langle 16_H \rangle (≡ \langle \bar{16}_H \rangle in the SUSY limit) for the gauge bosons and their superpartners of \{E, E^c\}, \ N, and \{Q, Q^c; U, U^c\}. The SO(10) gauge bosons associated with \{Q', Q'^c\} correspond to the so-called “X and Y” gauge bosons in SU(5). Hence, \langle 45_H \rangle and \langle 16_H \rangle determine the SU(5) and LR breaking scales, respectively.

Alternatively, one can employ the large representations, \{126_H, \bar{126}_H, and 210_H\}, instead of \{16_H, \bar{16}_H, and 45_H\} \[21\]. \ 126_H and \bar{126}_H break SO(10) to SU(5), while 210_H breaks SO(10) to SU(4)_c × SU(2)_L × SU(2)_R. In our discussion throughout this paper, \ 16_H (\bar{16}_H) and 45_H can be replaced by \ 126_H (\bar{126}_H) and 210_H, respectively.

Non-zero VEVs of both \langle 45_H \rangle and \langle 16_H \rangle eventually give the SM gauge symmetry at low energies. If \langle 45_H \rangle > \langle 16_H \rangle, SO(10) is broken first to LR at a higher energy scale and further broken to the SM gauge group at lower energy scales. On the other hand, if \langle 45_H \rangle < \langle 16_H \rangle, SO(10) is broken first to SU(5) at a higher energy scale and then eventually to the SM gauge group at lower energy scales. While the SU(5) breaking scale by \langle 45_H \rangle could be inferred from the renormalization group (RG) running effects of the three MSSM gauge couplings to be of \(3 \times 10^{16}\) GeV, the LR (or equivalently B – L) breaking scale by \langle 16_H \rangle may not be pinned down in principle: from the seesaw mechanism for the extremely light neutrinos, the LR breaking scale is just roughly estimated to be around \(10^{16}\) GeV. However, one should note that when the physical neutrino mass scale is theoretically estimated through the seesaw mechanism, the unknown Yukawa couplings associated with the RH neutrinos are involved. Moreover, the absolute neutrino masses can not be determined from the solar and atmospheric neutrino oscillations.
FIG. 1: Dominant diagram of the bino decay (a) and the gauge interaction between electrically charged superheavy LR gauginos and the MSSM lepton singlets (b).

Thus, if $\langle 45_H \rangle > \langle 16_H \rangle = \langle 16_H \rangle \neq 0$ ($\langle 16_H \rangle = \langle 16_H \rangle > \langle 45_H \rangle \neq 0$), the gauge bosons and gauginos of $\{Q', Q_c\}$ achieve heavier (lighter) masses than those of $\{E, E^c\}$ and $N$. The masses of the gauge sectors for $\{Q, Q^c; U, U^c\}$ would be given dominantly by the heavier masses in any cases, since both $\langle 45_H \rangle$ and $\{\langle 16_H \rangle, \langle \overline{16}_H \rangle\}$ contribute to their masses. Accordingly, the comparison of e.g. the gaugino masses of $\{Q', Q^c\}$ and $\{E, E^c\}$ (≡ $M_{Q'}$, $M_E$, respectively) could determine the hierarchy between $\langle 45_H \rangle$ and $\langle 16_H \rangle$, and so the SO(10) breaking pattern too.

III. LSP DECA Y IN SO(10)

If (1) R-parity is absolutely preserved and (2) $\chi$ is really the LSP, $\chi$ can never decay. We mildly relax these two conditions: by assuming a non-zero VEV of the superpartner of the (first family of) RH neutrino, $\tilde{\nu}^c_1$ (i.e. R-parity violation), or its mass lighter than the $\chi$’s mass, $m_\chi$ (i.e. $\tilde{\nu}_1$ LSP), $\chi$ can decay. By introducing a global symmetry, one can forbid its renormalizable Yukawa couplings to the MSSM fields. Then, $\tilde{\nu}^c_1$ can interact with the MSSM fields only through the superheavy gauge fields and gauginos of SO(10), since the (s)RH neutrino $\nu^c_1$ ($\tilde{\nu}^c_1$) is a neutral singlet under the SM gauge symmetry. Consequently, the decay of $\chi$ would be possible but quite suppressed. For instance, refer to the diagram of FIG. II(a). We will discuss how this diagram can be dominant for the $\chi$ decay.
A. The conditions for leptonic decay of $\chi$

Let us consider the interactions of the superheavy gauginos first. In TABLE I we list all the gauge interactions between the superheavy gauginos of SO(10) and two MSSM fields. They are, of course, the renormalizable operators. Since $\tilde{\nu}_i^c (i = 1, 2, 3)$ do not couple to $\tilde{Q}^c$ and $\tilde{Q}'$, the interactions by $\tilde{Q}^c$ and $\tilde{Q}'$ are not directly involved in the $\chi$ decay. As seen in TABLE II $\tilde{\nu}_i^c$ or $\nu_i^c$ couples to the superheavy SO(10) gauginos, $\{\tilde{E}, \tilde{E}^c\}$, $\tilde{N}$, $\{\tilde{Q}, \tilde{Q}^c\}$, and $\{\tilde{U}, \tilde{U}^c\}$.

According to PAMELA data [2], the branching ratio of the hadronic DM decay modes should not exceed 10%. To make the leptonic interactions, i.e. $\tilde{e}_i^{c*}\nu_i^c\tilde{E}^c$, $\tilde{\nu}_i^c\tilde{e}_i^c\tilde{\nu}^c$, and $\tilde{e}_i^{c*}\nu_i^c\tilde{N}$, $\tilde{e}_i^{c*}\nu_i^c\tilde{N}$ dominant over the other interactions in TABLE II we assume that

- The LR (or B – L) breaking scale should be lower than the SU(5) breaking scale, i.e. $\langle 16_H \rangle \ll \langle 45_H \rangle$. Then $M_{Q'}$, $M_Q$, $M_U$ (and also the masses of the superheavy triplet higgsinos contained in $10_h$) become much heavier than $M_E$ and $M_N$, and so most of hadronic decay modes of $\chi$ can be easily suppressed except those by $\tilde{E}^c$, $\tilde{E}$, and $\tilde{N}$ in TABLE II.

- The slepton $\tilde{e}_i^c$, which composes an SU(2)$_R$ doublet together with $\nu_i^c$, needs to be lighter than the squarks. Then the decay channels of $\chi$ by $\tilde{d}_i^{c*}\tilde{d}_i^c\tilde{E}^c$, $\tilde{\nu}_i^{c*}\nu_i^c\tilde{\nu}^c$, and $\tilde{u}_i^{c*}\tilde{u}_i^c\tilde{N}$, $\tilde{d}_i^{c*}\tilde{d}_i^c\tilde{N}$ become suppressed. We also require that $\chi$ and $\tilde{e}_i^c$ are much lighter than the charged MSSM.
Higgs. So the leptonic interactions, $\bar{\nu}_c^1 \nu_c^1 \bar{E}^c$, $\nu_c^1 \bar{E}^c$, and $\nu_c^1 \nu_c^1 \bar{N}$, $\bar{e}_c^1 e_c^1 \bar{N}$ can dominate over the others.

- At least one RH neutrino, i.e. the SU(2)$_L$ singlet neutrino $\nu_c^1$ (and its superpartner $\tilde{\nu}_c^1$) must be lighter than $\chi$ so that $\chi$ decays to charged leptons. It is because $\nu_c^i$ is always accompanied by $\tilde{\nu}_c^i$ in the effective operators leading to the leptonic decay of $\chi$, composed of $\bar{e}_c^i \nu_c^i \bar{E}^c$, $\nu_c^i \bar{E}^c$, and $\nu_c^i \nu_c^i \bar{N}$, $\bar{e}_c^i e_c^i \bar{N}$. If all the sneutrino masses are heavier than $\chi$, $\tilde{\nu}_c^i$ must develop a VEV for decay of $\chi$. Once $\nu_c^i$ is light enough, $\tilde{\nu}_c^i$ can achieve a VEV much easily.

To be consistent with PAMELA’s observations on high energy galactic positron excess[2], the DM mass should be around 300 – 400 GeV[13]. Thus, one can simply take the following values;

1. $\langle 16_H \rangle$ (or $\langle \tilde{\nu}_c^c \rangle$) $\ll \langle 45_H \rangle$. If $m_{\nu_c^i} > m_\chi$, then $\langle \tilde{\nu}_c^i \rangle \neq 0$.

2. squarks, charged Higgs, higgsinos and other typical soft masses are of $\mathcal{O}(1)$ TeV.

3. $m_{\nu_c^i} \ll m_\chi \sim 300 – 400 \text{ GeV} \lesssim m_{\tilde{\nu}_c^i} \ll \mathcal{O}(1) \text{ TeV}$. Consequently, SO(10) is broken first to LR, which would be the effective gauge symmetry valid below the GUT scale. As seen from TABLE II the gauge interactions by the LR gauginos (and also gauge fields) preserve the baryon numbers. Even if the masses of the LR gauginos and gauge fields are relatively light, their gauge interactions don’t give rise to proton decay. We will show later that the decay channels of $\chi$ through the mediation of the superheavy gauge fields are relatively suppressed.

### B. Seesaw mechanism

Although one RH neutrino is light enough, the seesaw mechanism for obtaining the three extremely light physical neutrinos still may work. Let us consider the following superpotential;

$$W_\nu = y_{ij}^{(\nu)} l_i h_u \nu_c^j (j \neq 1) + \frac{1}{2} M_{ij} \nu_c^i \nu_c^j (i, j \neq 1),$$

(2)

where the Majorana mass term of $\nu_c^i$ could be generated from the non-renormalizable superpotential $\langle 16_H \rangle \langle 16_H \rangle 16_{ij}/M_P$ ($i,j \neq 1$). Thus, $M_{ij}$ (or $\langle h_u \rangle$) could be determined, if the LR breaking scale by $\langle 16_H \rangle$ is known. In this superpotential, we note that one RH neutrino $\nu_c^1$ does not couple to the MSSM lepton doublets and Higgs. For instance, by assigning an
exotic U(1) R-charge to $\nu_1^c$, one can forbid its Yukawa couplings to the MSSM superfields. Thus, $\nu_1^c$ would be decoupled from the other MSSM fields, were it not for the heavy gauge fields and gauginos of the SO(10) SUSY GUT.

Taking into account only Eq. (2), one neutrino remains massless. The two heavy Majorana mass terms of $\nu_2^c$ and $\nu_3^c$ are sufficient for the other two neutrinos to achieve extremely small physical masses through the constrained seesaw mechanism [22]:

$$m_\nu = m_\nu^T = -\begin{pmatrix} 0 & v_{12} & v_{13} \\ 0 & v_{22} & v_{23} \\ 0 & v_{32} & v_{33} \end{pmatrix} \begin{pmatrix} 0 & 0 & 0 \\ 0 & M_{22}^{-1} & M_{23}^{-1} \\ 0 & M_{23}^{-1} & M_{33}^{-1} \end{pmatrix} \begin{pmatrix} v_{12} & v_{22} & v_{32} \\ v_{13} & v_{23} & v_{33} \end{pmatrix},$$

where $v_{ij} \equiv y^{(\nu)}_{ij}\langle h_u \rangle$, and $M_{ij}^{-1}$ denotes the inverse matrix of $M_{ij}$. One of the eigenvalues of $m_\nu$ is zero and the other two are of order $v^2/M$. This mechanism provides mixings of order $v/M$ between the three left-handed and two RH neutrinos. Through the diagonalization of the mass matrix in Eq. (3), the three left-handed neutrinos from the lepton doublet $l_1$, $l_2$, and $l_3$ can be maximally mixed, whereas the mixing of the RH neutrinos is only between $\nu_2^c$ and $\nu_3^c$. A complex phase in $y^{(\nu)}_{ij}$ could make leptogenesis possible [22].

C. Heavy gauginos’ masses

The gauge interactions between the gauginos and an SU(2)$_R$ lepton doublet $(2_1)$ in the LR model is described by

$$\mathcal{L} \supset -\frac{1}{2} \left( \tilde{e}^c \tilde{\nu}^c \right) \left[ g\tilde{N}_R + g'\tilde{N}_{BL} \right. \sqrt{2}g\tilde{E} \left. -g\tilde{N}_R + g'\tilde{N}_{BL} \right] \left( \begin{array}{c} e^c \\ \nu^c \end{array} \right) + \text{h.c.},$$

where $\{\tilde{N}_R, \tilde{E}, \tilde{E}^c\}$ and $\tilde{N}_{BL}$ are the superpartners of the SU(2)$_R$ and U(1)$_{B-L}$ gauge fields, respectively. $(-g\tilde{N}_R + g'\tilde{N}_{BL})/\sqrt{g^2 + g'^2}$ is identified with “$\tilde{N}$” discussed above. Hence, its orthogonal component $(g'\tilde{N}_R + g\tilde{N}_{BL})/\sqrt{g^2 + g'^2}$ corresponds to the bino of the MSSM. The hypercharge of the MSSM is defined by

$$\frac{Y}{2} = \pm \frac{1}{2}\sigma^3 + \frac{B - L}{2},$$

where $+(-)$ for 2 ($\overline{2}$). It is straightforward to write down the interaction between the LR gauginos and $\overline{2}_{-1}$. When the LR model embedded in the SO(10) GUT, the LR and B – L
gauge couplings, $g$ and $g'$ can be expressed in terms of the SO(10) gauge coupling,

$$g = \sqrt{\frac{2}{3}} g' = g_{10}. \quad (6)$$

By introducing a pair of SU(2)$_R$ doublet Higgs [or 16$_H$ and 16$^{-}_H$ in SO(10)],

$$\begin{align*}
2_1 &= \left( \begin{array}{c} e^c_H \\ \nu^c_H \end{array} \right) \subset 16_H, \quad \text{and} \quad \overline{2}_- = \left( \begin{array}{c} e_H \\ -\nu^c_H \end{array} \right) \subset 16^{-}_H,
\end{align*} \quad (7)$$

and, for instance, the superpotential

$$W = S \left( 2_1 \overline{2}_- - M_{LR}^2 \right) \subset S \left( 16_H 16^{-}_H - M_{LR}^2 \right), \quad (8)$$

one can break LR to the MSSM gauge group. Here, $S$ is a singlet superfield. By non-vanishing VEVs along the neutrino direction and the “D-flat” condition, $\langle \tilde{\nu}^c_H \rangle = \langle \tilde{\nu}^{-}_H \rangle = v/\sqrt{2}$, $\{e^c_H, \tilde{E}\}$ and $\{e_H, \tilde{E}^c\}$ obtain the same Dirac masses, and also the neutral gaugino $\tilde{N}$ and $(\nu^c_H - \nu^{-}_H)/\sqrt{2}$ ($\equiv \nu^c$) achieve a mass:

$$-\mathcal{L}_{mass} = M_E \left( e^c_H \tilde{E} + e_H \tilde{E}^c \right) + M_N \tilde{N} \nu^c$$

$$+ m_{3/2} \tilde{E} \tilde{E}^c + \frac{1}{2} m'_{3/2} \tilde{N}^2 + m''_{3/2} \left( e^c_H e_H + \frac{1}{2} \nu^c \nu^c \right) + \text{h.c.}, \quad (9)$$

where $M_E \equiv v_{g10}/2$ and $M_N \equiv v \sqrt{g^2 + g'^2}/2 = v_{g10} \sqrt{5}/8 = M_E \sqrt{5}/2$. We note here that $M_N$ is heavier than $M_E$. The other combination $(\nu^c_H + \nu^{-}_H)/\sqrt{2}$ ($\equiv \nu^c$) and $S$ get a mass from the superpotential Eq. (8) at the SUSY minimum. The second line of Eq. (9) contains the soft mass terms. Since $S$ can develop a VEV of order the gravitino mass $m_{3/2}$ due to the “A-term” corresponding to $W$ of Eq. (8), the last two mass terms of Eq. (9) [\subset \langle \tilde{S} \rangle (e^c_H e_H - \nu^c_H \nu^c_H)] are induced. We rewrite Eq. (9) in terms of the four component spinors as follows:

$$-\mathcal{L}_{mass} = \left( \lambda^{-}_R \psi^{-}_R \right) \left[ \begin{array}{c} m_{3/2} \\ M_E \\ m''_{3/2} \end{array} \right] \left( \begin{array}{c} \lambda^0_L \\ \lambda^0_R \psi^0_R \\ \psi^0_L \end{array} \right) + \frac{1}{2} \left( \lambda^0_R \psi^0_R \right) \left[ \begin{array}{c} m_{3/2} \\ M_N \\ m''_{3/2} \end{array} \right] \left( \begin{array}{c} \lambda^0_L \\ \psi^0_L \end{array} \right) + \text{h.c.} \quad (10)$$

where $\lambda^{-}(0)$ and $\psi^{-}(0)$ are the Dirac (Majorana) spinors constructed with the two components’ Weyl spinors for the gauginos and higgsinos:

$$\lambda^{-} = \left( \begin{array}{c} \tilde{E} \\ E^c \end{array} \right), \quad \psi^{-} = \left( \begin{array}{c} e_H \\ e^c_H \end{array} \right), \quad \text{and} \quad \lambda^0 = \left( \begin{array}{c} \tilde{N} \\ N \end{array} \right), \quad \psi^0 = \left( \begin{array}{c} \nu^c \\ \nu^{-} \end{array} \right), \quad (11)$$
where the “bar” denotes the complex conjugates of the fermionic fields. \( \lambda^+ \) and \( \psi^+ \) are respectively given by \((\lambda^-)^C\) and \((\psi^-)^C\), and \( \lambda^0 \) and \( \psi^0 \) satisfy \((\lambda^0)^C = \lambda^0 \) and \((\psi^0)^C = \psi^0 \). The mass eigenstates and their eigenvalues turn out to be
\[
\begin{pmatrix}
\Lambda_1^{-0} \\
\Lambda_2^{-0}
\end{pmatrix}_L = \frac{1}{\sqrt{2}} \begin{pmatrix}
1 - \epsilon & - (1 + \epsilon) \\
1 + \epsilon & 1 - \epsilon
\end{pmatrix} \begin{pmatrix}
\lambda_1^{-0} \\
\psi_1^{-0}
\end{pmatrix}_L, \quad \text{and}
\]
\[
M_{1,2}^{(-)} = \mp M_E + \frac{1}{2} \left[ m_{3/2} + m_{3/2}' \right], \quad M_{1,2}^{(0)} = \mp M_N + \frac{1}{2} \left[ m_{3/2}' + m_{3/2}'' \right],
\]
where \( \epsilon \equiv [m_{3/2}' - m_{3/2}] / (4M_{E,N}) \ll 1 \).

D. Heavy gauginos’ propagations

From Eq. (11), the charged interactions read as
\[
- \mathcal{L}_{c.c.} = \frac{g}{\sqrt{2}} \left[ \bar{e}^{cs}_i \nu^c_i \bar{E}^c + \bar{\nu}^{cs}_i e^c_i \bar{E} + \text{h.c.} \right] = \frac{g_{10}}{\sqrt{2}} \left[ \bar{e}^{cs}_i \bar{\lambda}^c P_L(\nu_{Di})^C + \bar{\nu}^{cs}_i \bar{e}^{c}_i P_L \lambda^- + \text{h.c.} \right],
\]
where \( P_L \) stands for the projection operator. \( \nu_{Di} \) and \( \bar{e}^{cs}_{Di} \) are Dirac spinors defined as
\[
\nu_{Di} = \begin{pmatrix} \nu \\ \bar{p}^c \end{pmatrix}_i \quad \text{and} \quad \bar{e}^{cs}_{Di} = \begin{pmatrix} \bar{e}^- \\ \bar{e}^c \end{pmatrix}_i.
\]
By contraction of \( \lambda^- \) and \( \bar{\lambda}^- \) in Eq. (14), therefore, the effective operator leading to \( \bar{e}^{cs}_i \rightarrow \bar{e}^-_i + \nu_1 + \bar{\nu}^c_1 \) is induced. See the diagram of Fig. 1(b). \( \lambda^- \) is decomposed to the two mass eigenstates \( \Lambda_1^- \) and \( \Lambda_2^- \), as shown in Eq. (12). With Eqs. (12) and (13), the amplitude suppression coming from the superheavy gaugino’s propagator \( \langle T\lambda^- \bar{\lambda}^- \rangle \) is estimated as
\[
i \left( \frac{1 - \epsilon}{\sqrt{2}} \right)^2 P_L \bar{\nu}^c + M_{1} \bar{P}_L + i \left( \frac{1 + \epsilon}{\sqrt{2}} \right)^2 P_L \bar{\nu}^c + M_{2} \bar{P}_L \approx i \frac{m_{3/2}^2}{M_{E}^2} P_L
\]
at low energies. Thus, the decay, \( \bar{e}^{cs}_1 \rightarrow \bar{e}^-_1 + \nu_1 + \bar{\nu}^c_1 \) is extremely suppressed, but still possible if it is kinematically allowed.

Eq. (11) includes also the neutral interactions of the SU(2)_L lepton singlets with \( \tilde{N} \) and the bino. One can extract the part interacting only with \( \tilde{N} \):
\[
- \mathcal{L}_{n.c.} = \frac{\sqrt{g^2 + g'^2}}{2} \bar{\nu}^{cs}_i \nu^c_i \bar{N} + \frac{g^2 - g'^2}{2\sqrt{g^2 + g'^2}} \bar{e}^{cs}_i e^c_i \bar{N} + \text{h.c.}
\]
\[
= \frac{g_{10}}{\sqrt{2}} \left[ \sqrt{5} \bar{\nu}^{cs}_i \bar{\lambda}^0 P_L(\nu_{Di})^C + \frac{1}{\sqrt{20}} \bar{e}^{cs}_i \bar{e}^{c}_{Di} P_L \lambda^0 + \text{h.c.} \right].
\]
They are actually reminiscent of the $Z$ boson interactions in the SM. By contracting $\lambda^0$ and $\overline{\lambda^0}$, the decay $\tilde{e}_1^c \rightarrow e_1^- + \nu_1 + \tilde{\nu}_1^c$ is possible. See FIG.2(a). However, since $M_N^2$ is $\frac{5}{2}$ times heavier than $M_E^2$ as shown from Eq. 9, and the effective coupling is $\frac{\sqrt{5}}{2} \times \frac{1}{\sqrt{20}} = \frac{1}{4}$ times smaller than that of the charged interaction case, the amplitude mediated by $\lambda^0$ is just $\frac{1}{10}$ of that by $\lambda^-$.

As seen in TABLE 1, the MSSM Higgs and higgsinos also couple to $\tilde{E}_1^c$, $\tilde{E}$ or $\tilde{N}$. Since the MSSM charged Higgs and higgsinos are assumed to be much heavier than $\tilde{e}_1^c$ and $\chi$, the decay channels through them are quite suppressed or kinematically forbidden.

So far we did not discuss the case in which $\chi$ decays through the mediation of the superheavy gauge bosons. The potentially dominant diagram is displayed in FIG.2(b). $\tilde{e}_1^c$ is coupled to $\chi$ and $e_1^c$. The scalar-scalar-gauge boson vertex is basically a derivative coupling. Accordingly, this diagram is suppressed compared to FIG.1(b), only if the bino is much lighter than the soft mass of $\{\tilde{E}, \tilde{E}^c\}$. As presented above, in this paper, we assume that $m_\chi \sim 300 - 400$ GeV and the soft mass of $\{\tilde{E}, \tilde{E}^c\}$ is of $O(1)$ TeV.

E. LSP decay rate and the seesaw scale

Now let us estimate the decay rate of FIG.1(a), which is the dominant decay channel, and determine the LR breaking scale such that it is consistent with PAMELA data. Indeed, if $m_{\tilde{\nu}_1^c} < m_\chi$, a non-zero VEV of $\tilde{\nu}_1^c$ is not essential: $\chi$ can decay to the four light particles, $e^\pm$, $\nu_1^c$, and $\tilde{\nu}_1^c$. However, just for simplicity, we will assume that a non-zero VEV of $\tilde{\nu}_1^c$ is
developed. For instance, let us consider the following terms in the superpotential;

\[ W \supset \frac{1}{M_P} \langle \mathbf{16}_H \rangle \mathbf{16}_1 \Sigma^2 + \kappa \Sigma^3, \]

where \( M_P = 2.4 \times 10^{18} \text{ GeV} \) and \( \kappa \) is a dimensionless coupling constant. \( \Sigma \) is an SO(10) singlet. We assign e.g. the U(1) R-charge of \( 2/3 \) to \( \mathbf{16}_1 \) and \( \Sigma \), and 0 to \( \mathbf{16}_H \). The scale of \( \langle \mathbf{16}_H \rangle = \langle \mathbf{45}_H \rangle = M_E/\sqrt{2} g_{10} \) can be determined such that it is consistent with PAMELA data. The soft mass term of \( \Sigma \) and the A-term corresponding to \( \kappa \Sigma^3 \) in the scalar potential permit a VEV \( \langle \tilde{\Sigma} \rangle \sim m_{3/2}/\kappa \). Then, the scalar potential generates a linear term of \( \tilde{\nu}_i \) coming from the A-term corresponding to the first term of Eq. (18), \( V \supset m_{3/2} (\langle \mathbf{16}_H \rangle/\kappa^2 M_P) \tilde{\nu}_i \). The linear term and the soft mass term of \( \tilde{\nu}_i \) in the scalar potential can induce a non-zero VEV of \( \langle \tilde{\nu}_1 \rangle \sim m_{3/2}/\kappa^2 \times M_E \).

Thus, the decay rate of \( \chi \) in FIG.1-(a) can be estimated:

\[ \Gamma_{\chi} = \frac{\alpha_{10}^2 \alpha_Y m_\chi^5}{96 M_E^4} \left( \frac{m_{3/2} \langle \tilde{\nu}_1 \rangle}{m_{\tilde{e}_1}} \right)^2 \sim \frac{\alpha_{10}^2 \alpha_Y m_\chi^5}{96 M_E^2 M_P^2} \left( \frac{m_{3/2}}{\kappa m_{\tilde{e}_i}} \right)^4 \sim 10^{-26} \text{ sec}^{-1}, \]

where \( \alpha_{10} (\equiv g_{10}^2/4\pi) \) and \( \alpha_Y \equiv g_Y^2/4\pi = (3/5) \times g_1^2/4\pi \), where \( g_1 \) is the SO(10) normalized gauge coupling of \( g_Y \) are approximately 1/24 and 1/100, respectively. Here, we ignore the RG correction to \( \alpha_{10} \). 300 – 400 GeV fermionic DM decaying to \( e^\pm \) and a light neutral particle can fit the PAMELA data \[13\]. For \( m_\chi \approx 300 – 400 \text{ GeV} \), \( \langle \mathbf{16}_H \rangle \sim 10, M_E \) or \( \langle \mathbf{16}_H \rangle \) is estimated to be of order \( 10^{14} \text{ GeV} \). This is consistent with the assumption \( \langle \mathbf{16}_H \rangle \ll \langle \mathbf{45}_H \rangle \sim 10^{16} \text{ GeV} \). Therefore, the masses of the other two RH neutrinos, which do not contribute to the process of FIG.1-(a), are around \( 10^{10} \text{ GeV} \) or smaller in this case: \( W \supset y_{ij} (\langle \mathbf{16}_H \rangle/\langle \mathbf{16}_H \rangle/M_P) \mathbf{16}_i \mathbf{16}_j (i, j \neq 1) \supset y_{ij} (10^{10} \text{ GeV}) \times \nu_i \nu_j^\dagger (i, j \neq 1) \). So the Yukawa couplings of the Dirac neutrinos should be a bit small (\( \sim 10^{-2} \)).

If \( m_\chi \approx 3.5 \text{ TeV} \) and the model is slightly modified such that \( \chi \) decays dominantly to \( \mu^\pm, \nu_2^\pm \) rather than to \( e^\pm, \nu_1^\pm \), which is straightforward, the Fermi-LAT’s data as well as the PAMELA’s can be also explained \[13\]. In this case, \( M_E \) or \( \langle \mathbf{16}_H \rangle \) should become somewhat heavier (\( \sim 10^{15} \text{ GeV} \)), and the seesaw scale should be replaced by \( 10^{12} \text{ GeV} \). However, the motivation of introducing SUSY to resolve the gauge hierarchy problem in the SM would become more or less spoiled.
IV. CONCLUSIONS

In this paper, we have shown that the bino-like LSP in the MSSM can decay through the SO(10) gauge interactions, if a RH neutrino is light enough \( m_{\nu_1} \lesssim m_\chi \) and its superpartner develops a VEV \( \langle \tilde{\nu}_1 \rangle \neq 0 \). The Yukawa couplings between the RH (s)neutrino and the MSSM fields can be suppressed by a global symmetry such as the U(1) R-symmetry. It gives rise to an extremely small R-parity violation very naturally. If the LR breaking scale or the seesaw scale is low enough compared to the GUT scale (i.e. \( \langle 16_H \rangle \ll \langle 45_H \rangle \sim 10^{16} \) GeV), and squarks, the MSSM charged Higgs, higgsinos, and other typical soft masses are relatively heavier (\( \sim \mathcal{O}(1) \) TeV) than the slepton, the recently reported PAMELA’s high energy galactic positrons can be explained through the leptonic decay of the bino-like LSP in the framework of the SO(10) SUSY GUT. Particularly, we assumed the quite mild hierarchies for the (s)lepton mass parameters; \( m_{\nu_1} \ll m_\chi \sim 300–400 \) GeV \( \lesssim m_{\tilde{\nu}_1} \ll \mathcal{O}(1) \) TeV. In the benchmark model, \( \langle 16_H \rangle \sim \mathcal{O}(10^{14}) \) GeV, and the two RH neutrino masses turned out to be of order \( 10^{10} \) GeV or smaller. Even if one RH (s)neutrino is almost decoupled from the interactions of the MSSM, the extremely light three physical neutrinos and their oscillations still can be achieved through the seesaw mechanism.

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