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JHEP04(2009)044

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R-parity preserving super-WIMP decays

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ABSTRACT: We point out that when the decay of one electroweak scale super-WIMP state to another occurs at second order in a super-weak coupling constant, this can naturally lead to decay lifetimes that are much larger than the age of the Universe, and create observable consequences for the indirect detection of dark matter. We demonstrate this in a supersymmetric model with Dirac neutrinos, where the right-handed scalar neutrinos are the lightest and next-to-lightest supersymmetric partners. We show that this model produces a super-WIMP decay rate scaling as $m_\nu^4/(\text{weak scale})^3$, and may significantly enhance the fraction of energetic electrons and positrons over anti-protons in the decay products. Such a signature is consistent with the observations recently reported by the PAMELA experiment.

KEYWORDS: Beyond Standard Model, Supersymmetric Effective Theories, Cosmology of Theories beyond the SM

ARXIV EPRINT: [0812.0432](https://arxiv.org/abs/0812.0432)

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1 Introduction

The detection of non-gravitational cold dark matter (DM) interactions is an important goal that drives some aspects of modern particle physics, cosmology and astrophysics. The well-measured abundance of dark matter [1] does not specify the origin or specific nature of DM. In particle physics, when the DM particle masses are comparable to the electroweak scale ($m_\chi \sim v$) there are two distinct (well-motivated) possibilities.¹ The first one, is the weakly interacting massive particle (WIMP) framework, in which the initially thermalized abundance of WIMPs is reduced via their weak annihilation, starting at the particles freeze out temperatures $T_f \sim 0.05 m_\chi$ [2, 3]. The second possibility (the super-WIMP) postulates that the DM particles have an interaction rate *much* weaker than the weak interactions. In this case, one assumes that the initial cosmological abundance of such particles is small, and that the rate for their direct production from thermal SM states, controlled by the square of the super-WIMP coupling y_{SW} , remains small relative to the Hubble expansion rate throughout the entire history of the Universe. The ratio of the super-WIMP thermalization rate to the Hubble rate at the weak scale, $\Gamma_{\text{th}}/H \sim 10^{-4} y_{\text{SW}}^2 \times (M_{\text{Pl}}/\text{weak scale})$, gives a crude estimate for the resulting number densities of electroweak scale super-WIMPs weighted by entropy. Observations require this ratio be $\lesssim 10^{-12}$, which sets the benchmark value for the super-WIMP coupling, $y_{\text{SW}} \lesssim 10^{-12}$. The stability of electroweak scale super-WIMPs would have to be ensured by some parity in the dark sector.

While the WIMP framework gives hope to the goal of direct detection of DM [4–6] (see however [7]), electroweak scale super-WIMPs would necessarily have a very tiny interaction strength that generally cannot lead to direct detection. Conversely, indirect signatures of DM, such as energetic gamma-rays, positrons and anti-protons, may be created by either

¹Of course, it is also possible that nature could correspond to the parameter space between two extremes of WIMPs and super-WIMPs and/or have multiple DM states of each type.

WIMPs or super-WIMPs, provided that the latter decay with a lifetime longer than the age of the Universe.

The recent claim by the PAMELA collaboration² [14] of an excess in the positron fraction above 10 GeV, is broadly in agreement with expectations for WIMP annihilation [15–21] or WIMP/super-WIMP decay [22, 23] leading to an excess positron flux above the background of secondary positron production. Even though an astrophysical origin of this signal (unrelated to dark matter) is not ruled out, there have been numerous attempts to link the PAMELA positron excess to dark matter. Various analyses of the positron flux [24–27] find that the PAMELA result can be fit to models of dark matter, provided that the annihilation and/or decay rate to positrons satisfy the following criteria:

$$\text{WIMP annihilation :} \quad \langle\sigma v\rangle_{e^+} \sim O(3 \times 10^{-24} \text{ cm}^3 \text{ s}^{-1}) \times \left(\frac{m_\chi}{500 \text{ GeV}}\right)^2 \quad (1.1)$$

$$\text{WIMP/super-WIMP decay :} \quad \Gamma_{e^+} \sim O(10^{-51} \text{ GeV}) \times \frac{m_\chi}{1 \text{ TeV}}, \quad (1.2)$$

where m_χ is the mass of decaying/annihilating particles. One should keep in mind that these estimates have uncertainties both due to their dependence on the modeling of the propagation of positrons in the galaxy (see [28, 29] for a discussion) and due to uncertainties in the local dark matter energy density. In addition, PAMELA has also reported the measured flux of anti-protons [30], which is well described by standard astrophysical production of anti-protons [31].

While the annihilation rate in eq. (1.1) of WIMPs inside our galaxy is naively in conflict with the abundance-derived rate of $\langle\sigma v\rangle_{T_f} \simeq 3 \times 10^{-26} \text{ cm}^3 \text{ s}^{-1}$, it has been argued that this may not necessarily be the case [32–34], as significant enhancement factors are possible due to resonance annihilation and/or v^{-1} enhancement from the long-range attraction in the dark matter sector.³ This can boost the galactic annihilation rate far above $\langle\sigma v\rangle_{T_f}$. As no excess in the anti-proton flux is reported, the positron excess suggests mostly leptonic channels of annihilation and small rates of decay to quark-antiquark pairs for the DM.

As a “boost factor” of size $10 - 10^3$ needed for the interpretation of the PAMELA results via WIMP annihilation may be accomplished in many ways through model-building, an explanation via the decays of WIMPs or super-WIMPs may seem more *ad hoc*. In order to obtain the decay width of 10^{-51} GeV in eq. (1.2), one would typically introduce a decay constant $y_{\text{decay}} \sim 10^{-25} - 10^{-23}$, such that $\Gamma_{\text{decay}} \propto y_{\text{decay}}^2 \times (\text{weak scale})/16\pi^2$. Such a coupling is twenty orders of magnitude smaller than a typical WIMP coupling and more than ten orders orders of magnitude smaller than the upper bound on the super-WIMP coupling. Although possible, such an option is not appealing, yet in most of the decay scenarios discussed in the literature, such a coupling is introduced by hand to fit the PAMELA data.

In this note, we consider the generic possibility when there are more than one super-WIMP states present in the DM sector, and the decay of one state to another is kinematically allowed. Such a decay can happen only in the second order of the super-weak

²As well as the more recent claim by the ATIC collaboration of an electron excess at higher energies [8] and other past hints of energetic electron/positron excess in the galaxy at high energies [9–13].

³Other recent works motivated by PAMELA data include [35–45]

coupling constant, and can be *naturally* suppressed down to the level given by eq. (1.2):

$$(\text{superWIMP})_1 \rightarrow (\text{superWIMP})_2 + \text{SM particles}, \quad \Gamma \sim O(y_{\text{SW}}^4). \quad (1.3)$$

The fourth order of the superweak coupling, $O(y_{\text{SW}}^4)$, is necessary a tiny number not far from 10^{-50} , which may lead to a natural realization of the $O(10^{-51} \text{ GeV})$ decay width.

We demonstrate this scenario realized in a supersymmetric Standard Model (MSSM) with exact R -parity conservation extended with right-handed (RH) neutrino superfields and Dirac (or nearly Dirac) neutrino masses. In this model one identifies the Yukawa couplings in the neutrino sector $y_\nu \sim 10^{-(12-13)}$ with y_{SW} [46–48]. Analyzing the model, we find that the excess of energetic positrons and electrons may naturally arise from the decay of one RH sneutrino species into another $\tilde{\nu}_R^1 \rightarrow \tilde{\nu}_R^2 \ell^+ \ell^-$, and that for certain domains of the parameter space the production of anti-protons is inhibited. This particular realization of the decaying super-WIMP scenario is the subject of this paper.

2 Naturally unstable SUSY DM candidates

The spectrum of supersymmetric particles depends on the nature of supersymmetry breaking generating the soft masses of the MSSM Lagrangian. Currently, there is no overwhelming reason to adhere to any particular scheme of SUSY breaking; the only objective restriction on the SUSY-breaking mass pattern comes from the resulting phenomenology. Thus, the charged lightest supersymmetric particle (LSP) is ruled out due to the desire for a natural DM candidate, and a nearly exact R -parity is required to keep the LSP stable if it is to be the DM. Many WIMP candidates such as left-handed sneutrinos are disfavored [49] due to direct detection constraints.

However, in order to accommodate the observed neutrino oscillations⁴ the SM should be supplemented with right handed neutrinos ν_R^i . The corresponding MSSM contains right-handed neutrino superfields that include both the fermionic right-handed neutrinos and their scalar partners, the right-handed sneutrinos $\tilde{\nu}_R^i$. Further, any choice of scale for the Majorana mass for the RH neutrinos is technically natural, which leaves the Yukawa couplings in the neutrino sector a free parameter that can be varied in a wide range $10^{-13} \lesssim y_\nu^i \lesssim 1$. If neutrino masses are Dirac (or nearly Dirac) then the Yukawa couplings are close to the lower end of this range, and the $\tilde{\nu}_R^i$ obtain the dominant contribution to their masses from the soft SUSY breaking mass terms in the MSSM; two of the $\tilde{\nu}_R^i$ masses, resulting mostly from these mass terms, could be the LSP and the NLSP.

Super-WIMP DM made up of $\tilde{\nu}_R^i$ was investigated recently in [46–48] and found to be a viable option.⁵ Indeed, $\tilde{\nu}_R^i$ are in some ways a very promising candidate for DM in the galactic halo: they can have the right relic abundance [46–48], and represents cold, neutral, colourless particles that do not interfere either with primordial nucleosynthesis or stellar evolution. Obviously, $\tilde{\nu}_R^i$ DM is consistent with limits on the self-interactions of DM and

⁴See [50] for a review of the experimental evidence of neutrino oscillation and [51] for a recent discussion of the theoretical formulation of neutrino oscillation.

⁵Decays of sneutrinos through R -parity violation was considered as a source of the PAMELA excess in [52].

leaves no detectable signal in direct detection searches [53, 54] as their interactions are suppressed by y_ν .

Despite these facts, one could question the phenomenological consequences of light Dirac-like neutrino masses due to the required smallness of the Yukawa coupling. The SM Yukawa couplings already have a hierarchy of 10^6 between the smallest and largest Yukawas. This hierarchy, although puzzling, does not lead to problems with technical naturalness as chiral symmetry is restored in the limit of vanishing Yukawa couplings. The Yukawa couplings for Dirac neutrinos are very small $\mathcal{O}(10^{-13})$, which leads to an even larger Yukawa coupling hierarchy but are likewise technically natural. Further, Dirac neutrinos may not have lepton number violation which disfavors some forms of leptogenesis⁶ as candidate theories for generating the observed matter-antimatter asymmetry of the Universe. However, many alternatives for baryogenesis exist such as electroweak MSSM baryogenesis [56], the Affleck-Dine mechanism [57], or an effective modification of the EW phase transition [58–60].

The benefits of $\tilde{\nu}_R^i$ as DM arguably outweighs the theoretical costs for the reason mentioned in the introduction. If the PAMELA data is to be explained through the decay of a particle, then the decay width must be smaller than the width corresponding to the age of the universe $\Gamma_U = \tau_U^{-1} \sim 10^{-42}$ GeV. Such small widths are challenging to naturally produce for weak scale WIMPs. On the other hand, in our scenario the small decay width is linked to another very small number, namely $y_\nu^4 \sim (10^{-13})^4$, which is an interesting possibility.⁷ One can immediately see that the decay widths are such that $\Gamma_\nu \ll \Gamma_U$. These decays would still be occurring in the galactic halo leading to a primary source of positrons that could be the excess observed by PAMELA, and possibly explain the ATIC data as well.

2.1 Dirac $\tilde{\nu}_R^i$ as LSP and NLSP

Following the initial stage of cosmological evolution, all superpartners of active SM species (charged under gauge groups) decay to $\tilde{\nu}_R^i$ on the time scales controlled by $\Gamma \propto y_\nu^2 \times$ (weak scale), see refs. [47, 48] for details. For our scenario, it is sufficient to assume that the population of the NLSP ($\tilde{\nu}_R^1$) is not parametrically small compared to the LSP ($\tilde{\nu}_R^2$) due to a very mild hierarchy of Yukawa couplings.

Consider the R -parity conserving MSSM supplemented with three Dirac $\tilde{\nu}_R^i$ with the superpotential

$$W = y_\nu^{ij} \epsilon_{\alpha\beta} \hat{H}_{u\alpha} \hat{L}_\beta^j \hat{\nu}_R^i - \mu \epsilon_{\alpha\beta} \hat{H}_{u\alpha} \hat{H}_{d\beta} + \dots \quad (2.1)$$

Here $\hat{H}_{u,\alpha} = (\hat{H}_u^+, \hat{H}_u^0)$ and $\hat{H}_{d,\alpha} = (\hat{H}_d^0, \hat{H}_d^-)$ are the up and down type Higgs chiral superfields and $\hat{L}_\beta = (\hat{\nu}_L, \hat{\ell}_L^-)$ is the lepton chiral superfield and $\epsilon_{12} = 1$. We have not written the terms in the superpotential containing the quark and charged lepton Yukawa couplings as they are not directly relevant to our discussion. The Yukawa interactions are

⁶Note that leptogenesis with Dirac neutrinos can occur in some models, see [55] for example.

⁷This observation was made in [46], where it was pointed out that the lifetime of the excited state of RH sneutrinos exceeds the lifetime of the Universe.

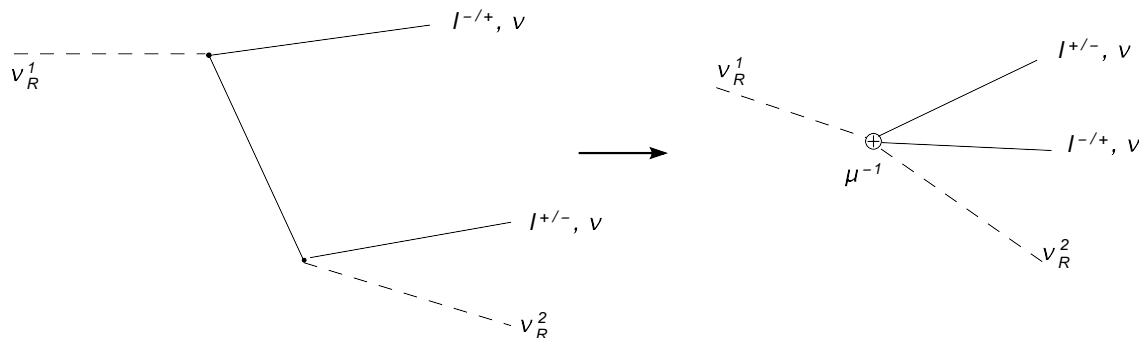


Figure 1. The tree level decay process for $\tilde{\nu}_R^1 \rightarrow \tilde{\nu}_R^2 \ell^\pm \ell^\mp$, $\tilde{\nu}_R^1 \rightarrow \tilde{\nu}_R^2 \nu \nu$ and the local operator approximation.

derived from the superpotential directly using

$$\mathcal{L} = \frac{1}{2} \sum_{i,j} \frac{\partial^2 W}{\partial \phi_i \partial \phi_j} \psi_i \psi_j + h.c. \quad (2.2)$$

Writing out the relevant component field terms in the resulting Lagrangian one has

$$\mathcal{L}_{\text{yuk}} = y_\nu^{ij} \left(\tilde{\nu}_R^i \nu^j \tilde{H}_u^0 - \tilde{\nu}_R^i \ell^j \tilde{H}_u^+ \right) - \mu \left(\tilde{H}_u^+ \tilde{H}_d^- - \tilde{H}_u^0 \tilde{H}_d^0 \right) + h.c. + \dots \quad (2.3)$$

When the lowest lying supersymmetric particles are $(\tilde{\nu}_R^i, \tilde{\nu}_R^j)$ the decay of one of these species into another can proceed via $\tilde{\nu}_R^1 \rightarrow \tilde{\nu}_R^2 \bar{\nu}^k \nu^l$ or $\tilde{\nu}_R^1 \rightarrow \tilde{\nu}_R^2 \bar{\ell}^k \ell^l$. As we insist on exact R -parity, decays of $\tilde{\nu}_R^i$ to purely SM particles is forbidden, while all other supersymmetric final states are kinematically inaccessible.

2.2 Leptonic $\tilde{\nu}_R^i$ decays

The leptonic decays $\tilde{\nu}_R^1 \rightarrow \tilde{\nu}_R^2 \ell^k \ell^l$ proceed through Higgsino exchange⁸ and the rate can be easily calculated as the function of m_1 , m_2 and μ , while masses of the SM leptons can be safely neglected. The amplitude is given by

$$i\mathcal{A} = i(y_\nu^{1j} y_\nu^{2i*}) \bar{L}^j \frac{\not{p}}{p^2 - \mu^2} L^i. \quad (2.4)$$

Here L is the left-handed lepton doublet, 1, 2 and j, l are flavour indices of RH sneutrinos and SM leptons respectively, and p is the momentum carried by the Higgsino. Since all mass parameters are essentially free apart from the $m_u > m_1 > m_2$ constraint, after quoting a general formula, we present the answer in several different limits. If the Higgsino mass μ is larger than the sneutrino mass scale, the Higgsino propagator can be contracted, as shown in figure 1, and the decay is mediated at leading order in p/μ by the effective Lagrangian,

$$\mathcal{L}_{\text{eff}} = \frac{y_\nu^{1j} y_\nu^{2i*}}{\mu^2} (\tilde{\nu}_R^{2*} \partial_\mu \tilde{\nu}_R^1) \bar{L}^j \gamma_\mu L^i, \quad (2.5)$$

⁸We disregard the issue of mixing between Higgsinos and gauginos as the mixing cannot introduce more than $O(1)$ corrections.

which is simply a product of sneutrino and left-handed lepton currents. Introducing effective couplings $y_1^2 = \sum_j |y_\nu^{1j}|^2$, $y_2^2 = \sum_j |y_\nu^{2j}|^2$ and the energy release parameter $\Delta = m_1 - m_2$ one can calculate the general decay width in a straightforward manner. One finds the result for the decay width given by

$$\frac{\Gamma_{l^+l^-}}{y_1^2 y_2^2} = -\frac{\Delta (24m_1^7 - 84\Delta m_1^6 + 104\Delta^2 m_1^5 - 50\Delta^3 m_1^4 - 8\Delta^4 m_1^3 + 20\Delta^5 m_1^2 - 8\Delta^6 m_1 + \Delta^7)}{6144m_1^3 \pi^3 \mu^4},$$

$$-\frac{12 m_1^4 \log \left[\frac{(m_1 - \Delta)^2}{m_1^2} \right] (m_1 - \Delta)^4}{6144 m_1^3 \pi^3 \mu^4}. \quad (2.6)$$

This general result obscures the physics somewhat, however, recall that in our scenario the sneutrinos are the LSP and the NLSP so necessarily the Higgsino is a larger mass scale. Let $\mu = n m_1$ where n is an order one number ($n > 1$) for this reason. Further, for this decay to occur $m_2 < m_1$ so let $m_2 = x m_1$ where $x < 1$, then expanding in small x one obtains

$$\Gamma_{l^+l^-} = \frac{m_1 y_1^2 y_2^2}{6144 n^4 \pi^3} (1 - 8x^2 - 24x^4 \log(x) + 8x^6 + \mathcal{O}(x^7)). \quad (2.7)$$

Thus we see that a good order of magnitude approximation for this decay width is $\Gamma_{l^+l^-} \simeq 10^{-6} y_1^2 y_2^2 m_1$ for the whole range of parameter space.

We also note that this decay width has a number of phenomenologically interesting limits. For example, considering Δ to be somewhat small on the scale of the sneutrino masses, we derive the total decay rate into charged leptons,

$$\Gamma_{l^+l^-}^a = \frac{y_1^2 y_2^2 \Delta^5}{480 \pi^3 \mu^4} \quad \text{for } \Delta \ll m_{1,2} \ll \mu. \quad (2.8)$$

This formula can be obtained directly from the neutron beta decay rate upon setting $m_e, g_A \rightarrow 0$, and identifying $G_F \cos \theta_c / \sqrt{2} \rightarrow 1/(4\mu^2)$. In another kinematic regime, when μ is very close to m_1 and the energy release remains small, the rate is given by

$$\Gamma_{l^+l^-}^b = \frac{y_1^2 y_2^2 \Delta^3}{192 \pi^3 m_1^2} \quad \text{for } \mu - m_1 \ll \Delta \ll m_{1,2}. \quad (2.9)$$

Finally, if the scale of m_2 is small and can be neglected, while μ remains close to m_1 we have a third limit,

$$\Gamma_{l^+l^-}^c = \frac{y_1^2 y_2^2 m_1}{1024 \pi^3} \quad \text{for } \mu - m_1, m_2 \ll m_1. \quad (2.10)$$

We do not require that any of the hierarchies between mass parameters m_1, m_2, μ used to determine the widths $\Gamma_{l^+l^-}^{a,b,c}$ is strictly realized in our scenario. We have determined $\Gamma_{l^+l^-}^{a,b,c}$ using these limits on m_1, m_2, μ to determine some phenomenologically interesting limits of the leptonic decay width. In what follows we will use the approximation $\Gamma_{l^+l^-} \simeq 10^{-6} y_1^2 y_2^2 m_1$ which is a good approximation due to the physics of our scenario, (ie a Higgsino heavier than the LSP and NLSP sneutrinos and $m_2 < m_1$).

2.3 Hadronic $\tilde{\nu}_R^i$ decays as a source of \bar{p} flux

Depending on kinematics, besides the pure leptonic decay channels, there could be other decay channels potentially leading to energetic anti-quarks and eventually to anti-protons. Quark-antiquark pairs cannot be produced in the sneutrino decay via the intermediate gauge and/or Higgs bosons. The relevant part of the scalar potential that could lead to such decays is

$$V = 2\mu y_\nu^{ij} \tilde{\nu}_R^i (\tilde{\nu}_L^j H_d^0 - \tilde{\ell}_L^j H_d^+) + \sum_k y_\nu^{ik} (y_\nu^{jk})^* H_u^0 (H_u^0)^* \tilde{\nu}_R^i \tilde{\nu}_R^j + h.c. + \dots \quad (2.11)$$

These interactions induce three- and four-body decays of RH neutrinos: $\tilde{\nu}_R^1 \rightarrow \tilde{\nu}_R^2 h(H)$, $\tilde{\nu}_R^1 \rightarrow \tilde{\nu}_R^2 W^+ W^- (H^+ H^-)$, $\tilde{\nu}_R^1 \rightarrow \tilde{\nu}_R^2 ZZ(AA)$, decays. To simplify our calculations, we shall assume that $\tan\beta = \langle H_u^0 \rangle / \langle H_d^0 \rangle$ is somewhat large, and the A, H, H^\pm Higgs bosons are heavy, so that only the decay to the lightest Higgs boson h and pairs of gauge bosons are kinematically possible.

We concentrate on the three-body decay $\tilde{\nu}_R^i \rightarrow \tilde{\nu}_R^j h$, which is driven by the last term in eq. (2.11) when one of the Higgs is given a vev. The resulting decay width is

$$\Gamma_h = |y_{12}^2|^2 \frac{v^2}{8\pi m_1^2} \times v_h \left(\Delta + \frac{m_h^2 - \Delta^2}{2m_1} \right). \quad (2.12)$$

In this formula, v_h is the velocity of the outgoing Higgs boson, $v = 246$ GeV is the SM Higgs vev, and y_{12}^2 is the sneutrino-flavor changing combination of the Yukawa couplings:

$$y_{12}^2 \equiv \sum_j y_\nu^{1j} (y_\nu^{2j})^*. \quad (2.13)$$

If $m_1 \sim v$, $v_h \sim 1$ and $y_{12}^2 \sim y_1 y_2$, the two-body decay rate would dominate over the leptonic rate by three orders of magnitude simply because of the phase space suppression of eqs. (2.8)–(2.10) relative to eq. (2.12) is quite significant. The Higgs boson produced in the decay fragments further to gauge bosons and heavy quarks and leptons. For SUSY models its typical mass precludes it from decaying directly to $W^+ W^-$ or ZZ , and $b\bar{b}$ decays are expected to dominate. Decays of this form is often encountered in neutralino annihilation, and the corresponding yield of antiprotons from a b -quark injection has been evaluated [61, 62]. The yield of antiprotons in the hadronization and subsequent decays of $b\bar{b}$ -pairs is quite significant, exceeding 10% per annihilation/decay event with a typical electroweak scale energy injection. Assuming an $\mathcal{O}(0.1 - 1)$ yield of antiprotons in the $\tilde{\nu}_R^1 \rightarrow \tilde{\nu}_R^2 h$ decay, the ratio of eq. (2.12) to a typical three-body rate gives the relative strengths of antiprotons and positrons at the injection:

$$\frac{\Phi_{\bar{p}}}{\Phi_{e^+}} \sim 10^3 \times \frac{|y_{12}^2|^2}{y_1^2 y_2^2} \times v_h. \quad (2.14)$$

A recent theoretical study of the antiproton fraction [31] limits the "boost factor" for the annihilation of dark matter to be less than 6(40) for the annihilation of 0.1(1) TeV particles. If we tune the model to fit PAMELA flux for positrons, this would translate into

a constraint on eq. (2.14) to be smaller than 0.1-1. Consequently, one would have to require a suppression of $v_h |y_{12}^2|^2 / (y_1^2 y_2^2)$ down to the level of $10^{-4} - 10^{-3}$. Such a suppression may come from the smallness of the flavor-changing coupling y_{12}^2 , or may originate kinematically, from small scale splitting, $\Delta < m_h$. In both cases one would be able to fit PAMELA results without overproducing antiprotons.

Note that most decays that are loop suppressed will also have the difference in flavour decays to quarks and thus protons and anti-protons compared to leptonic decays. This is due to the exact R parity of our scenario and the fact that only the LSP $\tilde{\nu}_R^2$ is kinematically accessible for the NLSP $\tilde{\nu}_R^1$ to decay too. Loop decays with the same flavour structure exist where the produced leptons form a loop that produces a Z/γ that subsequently decays to hadrons. The most problematic of these decays are suppressed by

$$\frac{\Phi_{\bar{p}}}{\Phi_{e^+}} \sim 10^3 \times \frac{g^2 g_V^2}{16\pi^2 (4 \cos^2(\theta_W))} \times BR(Z \rightarrow \text{hadrons}). \quad (2.15)$$

and the suppression is smaller than the required 0.1-1.

2.4 Numerology

Our main result, eqs. (2.8)–(2.10) can give an excellent fit to the PAMELA data in agreement with the requirement of eq. (1.2). Adopting $\Gamma_{l^+ l^-} = 10^{-6} \times y_1^2 y_2^2 \times m_1$ as an approximation for different kinematic regimes of the three-body decay rate, we conclude that the Yukawa couplings must satisfy the following relation:

$$(y_1 y_2)^2 \sim 10^{-52} \implies (y_1 y_2)^{1/2} \sim 1 \times 10^{-13}. \quad (2.16)$$

Are these values for the Yukawa couplings in agreement with Dirac neutrino Yukawa couplings suggested by the measurements of neutrino oscillations? The flavor basis in the RH sneutrino sector and the active neutrino sector in general do not coincide and a one-to-one connection between neutrino phenomenology and eq. (2.16) is not possible. The flavor-changing combination of Yukawa couplings in eq. (2.12), remains uncertain and is not necessarily related to the measured mixing angles. If further assumptions are made about the structure of the PMNS and the V_R mass matrices in a particular constrained MSSM then the the flavour structure we have determined can be related to the PMNS mass matrix, but this is beyond the scope of this work. No immediate connection between our determined flavour structure and the observable masses and mixing angles in the neutrino sector is possible without further experimental input, but we can check the consistency of eq. (2.16) with general expectations for the Yukawa couplings in the Dirac neutrino sector. Recall that Dirac neutrinos obtain their masses via $m_\nu = y_\nu \langle H_u^0 \rangle = y_\nu v \sin \beta / \sqrt{2} \simeq y_\nu \times 175 \text{ GeV}$, as we assume $\tan \beta$ to be large. To estimate the possible size of the Yukawa couplings we use the experimental determinations of the neutrino mass splitting. The difference in masses squared of the neutrinos have been measured by the KamLAND [63] and K2K [64] collaborations to be

$$[\delta m_\nu^2]_{\text{atm}} \simeq 2.8 \times 10^{-3} \text{ eV}^2, \quad [\delta m_\nu^2]_{\text{solar}} \simeq 7.9 \times 10^{-5} \text{ eV}^2 \quad (2.17)$$

and the absolute bound on the sum of neutrino masses from WMAP is given to be

$$\sum m_\nu^i = 0.67 \text{ eV at } 95\% \text{ CL.} \quad (2.18)$$

which coincides with the limit on the heaviest neutrino mass $m_\nu < 0.7 \text{ eV}$.

In the case of the standard hierarchy of neutrino masses the heaviest neutrino is given by $\sim \sqrt{(\delta m_\nu^2)_{\text{atm}}}$ and one finds that the largest Yukawa coupling is

$$y_\nu \simeq 3.0 \times 10^{-13} \left(\frac{m_\nu^2}{2.8 \times 10^{-3} \text{ eV}^2} \right)^{1/2}, \quad (2.19)$$

which is perfectly consistent with the requirement on $(y_1 y_2)^{1/2}$ in our scenario. Alternatively, for the small neutrino mass splitting scenarios one could have Yukawa couplings close to the value suggested by possible given cosmological bounds, $y_\nu \simeq 10^{-12}$. In either case we see that the Yukawa couplings can realize our super-WIMP decay scenario that produces the PAMELA, and possibly ATIC, signals for electroweak scale DM masses.

3 Discussion and conclusions

We have presented an observation that the decay rates of super-WIMPs with the rate about ten orders of magnitude slower than the Hubble rate may originate from the transition of one super-WIMP state into another in the *second* order in the super-WIMP coupling so that $\Gamma \sim y_{\text{SW}}^4 \times (\text{weak scale})$. The resulting decay rates are just becoming possible to probe through indirect detection of dark matter via its decay products, such as electrons, positrons and antiprotons.

Perhaps the most interesting aspect of our model is that super-WIMP physics is extremely sensitive to the superpartner mass spectrum due to the existence of a very small coupling y_{SW} . If, for example, one has a mass spectrum such that $m_{\tilde{\nu}_R^2} < m_{\text{Higgsino}} < m_{\tilde{\nu}_R^1}$, the sequential decays $\tilde{\nu}_R^1 \rightarrow \text{Higgsino} \rightarrow \tilde{\nu}_R^2$ will happen with $\Gamma \propto y_\nu^2 \times (\text{weak scale}) \propto 10 \text{ Hz}$, which is almost instantaneous on the scale of eq. (1.2). A simple modification of the spectrum to $m_{\tilde{\nu}_R^1} < m_{\tilde{\nu}_R^2} < m_{\text{Higgsino}}$ was shown to lead to an enormous delay in the $\tilde{\nu}_R^1 \rightarrow \tilde{\nu}_R^2$ decay. This delay results in an overall drop in the decay rate by 25 orders of magnitude, and makes the heavier component of the RH sneutrinos only very weakly unstable. The numbers, given a huge disparities of different scales involved, work remarkably well, producing $\Gamma_{e^+} \sim m_\nu^4 / (\text{weak scale})^3$, which is in agreement with a putative explanation of PAMELA (and ATIC) signals by the decaying super-WIMPs. Moreover, for a moderate splitting between two RH sneutrino components and/or an accidental suppression of the different flavor transition, the hadronic decays of $\tilde{\nu}_R^1$ will be suppressed, which will in turn suppress the antiproton flux created by super-WIMPs. A kinematic suppression of the antiproton flux would imply a rather small energy splitting, $\Delta < m_h \sim O(100) \text{ GeV}$ between two super-WIMP states. Alternatively, if both PAMELA and ATIC are to be explained in the same way, one would have to choose Δ close to a TeV and assume a hierarchy $|y_{12}^2|^2 \ll y_1^2 y_2^2$ in order to suppress the antiproton flux.

One feature of our scenario may look somewhat unusual. We do not assume a complete degeneracy of RH sneutrino masses, and in fact require them being split by the energy

intervals comparable to their masses. This is not possible in the context of universal sfermion masses often employed in the SUSY literature, as any additional splitting induced by y_ν is minuscule. Therefore, one would have to imagine some additional theoretical mechanism to create such a splitting.

Although the example given in this note is quite natural, it is hardly unique. For example, one could envisage sequential R -parity preserving decays of other super-WIMPs featured in the SUSY literature, such as gravitinos, axinos, modul(inos) etc.

The decaying super-WIMP scenario as a tentative explanation of PAMELA and ATIC differs from the boosted WIMP annihilation in many ways. The most notable distinction is, of course, the early cosmology. The enhanced annihilation may lead to the nuclear-chemical consequences such as the overproduction of ${}^6\text{Li}$ [65], and to extra ionization and a diffuse gamma background [66] produced by the annihilating WIMPs. In contrast, the decaying DM scenario is immune to these potential problems. On the other hand, both annihilation and decay scenarios are subject to gamma ray and synchrotron emission constraints (for a recent paper see [67]) coming from the central region of the galaxy where the density of dark matter is significantly enhanced.

Acknowledgments

We would like to thank Adam Ritz, Brian Batell and Vivek Sharma for useful discussions. The work of MP was supported in part by NSERC, Canada. Research at the Perimeter Institute is also supported in part by the Government of Canada through NSERC and by the Province of Ontario through MEDT.

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