Two dark matter components in dark matter extension of the minimal supersymmetric standard model and the high energy positron spectrum in PAMELA/HEAT data

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We present a dark matter extension of the minimal supersymmetric standard model to give the recent trend of the high energy positron spectrum of the PAMELA/HEAT experiments. If the trend is caused indeed by dark matter, the minimal supersymmetric standard model needs to be extended. Here, we minimally extend the minimal supersymmetric standard model with one more dark matter component N together with a heavy lepton E and introduce the coupling $e_R E_R^c N_R$. This coupling naturally appears in the flipped SU(5) grand unification models. We also present the needed parameter ranges of these additional particles.

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The existence of dark matter (DM) at the 23% level of closure density [1] is largely accepted by the measurement of the velocity dispersion in the galaxy clusters, observations of the flat rotation curves of the velocities of stars, the x-ray emission, and the gravitational lensing effect from the galaxy cluster [2]. So, the identification of the cosmological DM is of prime importance in particle physics and cosmology. If indeed the DM particles of O(100) GeV mass with weak interaction strength are abundant in galaxies, high energy positrons, antiprotons, and gamma rays from DM annihilation have been predicted for a long time [2]. If the DM annihilation is confirmed, the *very weakly* interacting DM possibility [3] is ruled out [4].

The PAMELA experiment has already started to search for the DM signal. Their recent report on the high energy positron observation, above 10 GeV up to 80 GeV [5], has already spurred a great deal of attention [6,7]. In fact, the same trend has been noticed earlier in the balloon-borne HEAT experiment but with larger error bars [8]. The characteristic of the PAMELA/HEAT data between 10–80 GeV is a prominently rising positron flux, $e^+/(e^+ + e^-) \sim$ O(0.1). It is expected to be explained by DM annihilation. If the positron excess is caused indeed by DM and confirmed by independent observations such as PEBS balloon experiments [9] and the AMS-02 experiments [10], the implication is tantamount in that the most popular lightest neutralino χ (LNχ) DM scenario of the minimal supersymmetric standard model (MSSM) may be in jeopardy. Even though the possible astrophysical explanations have been presented in [11], here we focus on the particle physics explanation. On the other hand, we note that the same PAMELA data does not show any significant antiproton excess [5].

If cold dark matter (CDM) is composed of just one component Majorana fermion like LNχ, their fermionic nature severely constrains the annihilation channel. In general, the Fermi-Dirac statistics forces two identical Majorana particles in the s-wave state to be in the antiparallel spin state. Thus, at present, in most cases of the CDM annihilation ($v \sim 10^{-3}$), the angular momentum is initially zero. On the other hand, if the final state after the interaction is composed of two fermions $f\bar{f}$ much lighter than CDM, it can have the zero angular momentum only with the helicity flipping interaction, which is too small for the electron-positron pairs. Hence, the annihilation of neutralinos into c, b, and t quark pairs (and also W and Z boson pairs), if it is open, would dominate over all other channels including the channel to $e^\pm$ in the MSSM. In view of the conventional MSSM CDM scenario, therefore, the recently reported PAMELA data on the high energy positron excesses is quite embarrassing.

This leads us to consider a minimal extension of the MSSM so as to keep its most desirable property *"supersymmetry."* In the framework of spin-$\frac{1}{2}$ CDMs, therefore, we extend the MSSM minimally to include two CDM components.

In the MSSM, if the LNχ is bino denoted as χ, then the cosmologically favored bino density in the universe is possible in the coannihilation region [12]. In this paper we treat the LNχ as the bino for the sake of a concrete discussion. Our result is based on the assumption that LNχ contains the bino as a significant fraction. Otherwise, the ZZ and W^+W^- channels are open if kinematically allowed. Then, excess antiprotons are possible by gauge boson decays.

As commented above, two annihilating neutralinos of Fig. 1(a) are of zero angular momentum. If the final electron and positron, going out back to back, have the same helicity, then they can make up the angular momentum zero. One such possibility shown in Fig. 1(a) shows that it is highly suppressed because the bullet of Fig. 1(a) carrying an SU(2)$_W$ quantum number $f_e\langle H_d^0\rangle m_{3/2}$ is a high suppression factor. On the other hand, if the outgoing electron and positron carry opposite helicities, their spin is one and by emitting a photon the three particles final state can make up angular momentum zero. However, in this case there is a coupling suppression of order $\alpha_{\rm em}/\pi$. This possibility of high energy positron plus photon has

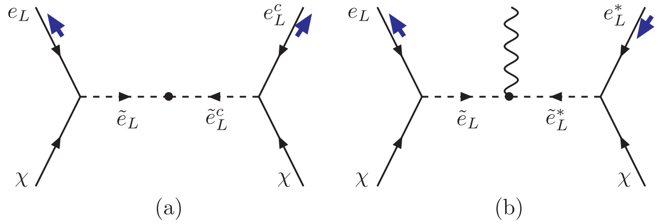

FIG. 1 (color online). The binolike neutralino annihilation. (a) The bullet carries an $SU(2)_W$ quantum number. (b) Here, the bullet can carry angular momentum 1. The helicities of electrons and positrons are shown by thick arrow lines.

been suggested in Ref. [6] where a huge boost factor from the DM halo clumpiness of order 10^4 should be assumed even in the fine-tuned parameter region. A scalar DM such as sneutrino LNχ will have the same fate as the bino LNχ in this regard.

Instead of the direct production of positrons, one may consider a heavy winolike neutralino which dominantly produces W^+W^- and ZZ so that subsequent decays provide positrons [13]. But this requires a large LNχ mass of a few TeV, and it spoils the naturalness for the small Higgs mass significantly. Moreover, the decays of W and Z possibly give some excess of the antiprotons also.

Without considering SUSY, the authors of [14] introduce vectorlike SU(2) multiplets ($I = 5$) which are annihilated into W^+W^- and ZZ. It avoids the antiproton constraint by raising the DM mass to O(10 TeV). However, the motivation of introducing the very heavy isospin multiplet is unclear.

Of course, a Dirac particle DM can be considered [15]. Also a spin-1 DM can overcome the binolike LNχ difficulty, but in the context of the minimal universal extra dimension model ("mUED") and the little Higgs model with T parity ("LHT") [7], it would give some excess of antiprotons. These considerations present a useful direction for constructing DM models with the rising high energy positron spectrum above 10 GeV.

Thus, we extend the MSSM so that the DM annihilation produces high energy positrons but not excess antiprotons, thus providing a possible explanation of the PAMELA/HEAT data. Accordingly, a special treatment of color singlet particles compared to quarks is necessary when extending the MSSM. Moreover, avoiding the "helicity suppression" by two identical Majorana fermions needs one more dark matter component. Thus, let us consider the extended MSSM with a neutral singlet superfield N, whose fermionic component contains the needed extra DM component ($N_{\rm DM}$MSSM). [Throughout this paper, we will use the same notations for the superfields and their fermion components, unless they give rise to serious confusion.] N is split into two Weyl spinors $\{N_R, N_L\}$. However, the coannihilation process, $\chi + N \to e^+ + e^-$, would kinematically allow also $\chi \to N + e^+ + e^-$ or $N \to \chi + e^+ + e^-$, unless the masses of χ and N are extremely fine-tuned ($|M_\chi - m_N| < 2m_e$). Thus, we also introduce an SU(2) *singlet* charged lepton E ($= \{E_R, E_L\}$), replacing e^- by E in $\chi + N \to e^+ + e^-$. This extra charged lepton makes χ and N stable in some region of the parameter space as we will discuss below.

This $N_{\rm DM}$MSSM seems to be very simple in the sense that new particles needed beyond the MSSM are *minimal*, just N and E. To be relevant in low energy physics, the SM singlet N needs to be light enough. Let us introduce the continuous R symmetry so that the additional DM component N remains light down to low energies. The weak hypercharge Y and R charge of the singlet fields are

$$\begin{array}{cccccc} \text{Superfields:} & e_R & N_R & N_R^c & E_R & E_R^c \\ Y: & -1 & 0 & 0 & -1 & +1 \\ R: & +1 & \frac{2}{3} & -\frac{2}{3} & -\frac{1}{3} & +\frac{1}{3} \end{array} \tag{1}$$

The R charges of the MSSM fields are as usual: the quark and lepton superfields carry 1 and the Higgs superfields carry 0. The R symmetry allows the superpotential

$$W = f e_R E_R^c N_R + h N_R^3, \tag{2}$$

where f and h are coupling constants, but the mass terms of N and E cannot be present in W. However, via the Giudice-Masiero mechanism [16], supergravity effects can generate the fermion masses if the F term of a singlet S is developed,

$$\int d^4\theta \left[\frac{S^*}{M_P} (\lambda E_R E_R^c + \lambda' N_R N_R^c) + \text{H.c.} \right]. \tag{3}$$

Similarly the MSSM μ term is also generated. These masses are assumed to be of order the gravitino mass $m_{3/2}$. N_R does not develop a VEV. As usual, with the separate lepton number conservation the process $\mu \to e\gamma$ is forbidden. Without the separate lepton number conservation, the $U(1)_R$ allows E_R^c couplings to μ_R and τ_R by couplings f' and f'', in which case we need $|f'| \lesssim 10^{-4}$. The f'' bound is much weaker.

Not introducing N couplings to quarks, we will not introduce excess antiprotons. From string compactification with the doublet-triplet splitting [17], the flipped SU(5) is best suited for this purpose [18]. In the flipped SU(5) whose gauge group is $SU(5) \times U(1)_X$, the hypercharge Y is given by a linear combination of the diagonal generator of SU(5) and the $U(1)_X$ generator, $Y = (Y_5 - X)/5$. Since N_R, e_R, and E_R^c remain as SU(5) singlets 1_0, 1_5, and 1_{-5}, respectively, the superpotential Eq. (2) is still invariant under the flipped SU(5). Quark fields in the flipped SU(5) appear in the representations **10** and $\bar{\mathbf{5}}$ and hence cannot couple to our DM candidate N_R which is a flipped SU(5) singlet. This is the reason that our model does not give excess antiprotons.

The mass terms induced by SUSY breaking can be written as an effective superpotential,

$$W \supset m_{3/2} E_R E_R^c + m'_{3/2} N_R N_R^c, \tag{4}$$

and soft SUSY breaking A and B terms are

$$\mathcal{L} \supset m_{3/2}[a_f \tilde{e}_R \tilde{E}_R^c \tilde{N}_R + a_h \tilde{N}_R^3] + m_{3/2}^2[b\tilde{E}_R\tilde{E}_R^c + b'\tilde{N}_R\tilde{N}_R^c] + \text{H.c.}, \tag{5}$$

where a and b denote dimensionless couplings. Equation (5) violates the continuous R symmetry, leading to a discrete Z_6 symmetry. The R parity (Z_2) is a part of this Z_6 symmetry. For the fields in Eq. (1), the (R parity, Z_6) charges are

$$\begin{gathered} \tilde{e}_R(-,3), \quad \tilde{N}_R(+,2), \quad \tilde{N}_R^c(+,4), \quad \tilde{E}_R(-,5), \\ \tilde{E}_R^c(-,1), \quad e_R(+,0), \quad N_R(-,5), \quad N_R^c(-,1), \\ E_R(+,2), \quad E_R^c(+,4). \end{gathered} \tag{6}$$

The SM fields are given $Z_6 = 0$ and their superpartners are given $Z_6 = 3$, which is just the R parity for this MSSM subset field in our $N_{\rm DM}$MSSM. The proton longevity is protected by the $U(1)_R$. The $U(1)_R$ forbids the dimension 5 operators $q_R^c q_R^c q_R^c l_R^c$ and $u_R u_R d_R e_R$, which are allowed by the conventional R parity alone, as well as the R parity violating the dimension 4 operators. If SUSY breaking induces them in the $N_{\rm DM}$MSSM, they must be highly suppressed since there are no simple diagrams for them. Because of the R- parity of N, N cannot be a candidate for the singlet heavy neutrino of the seesaw mechanism.

We intend to introduce two stable particles, one the LNχ and the other the lightest Z_6 matter particle (LMP). For the superfields $\tilde{N}_R + N_R\theta$ and $\tilde{E}_R + E_R\theta$, we assume $M_{\tilde{N}} > m_N$, $M_{\tilde{E}} > m_E$, and $m_E > m_N$. We assume that N is lighter than $\tilde{e}_R$ and $\tilde{E}_R$, and hence N is taken as the LMP. The LMP N carries $Z_6 = 5$ which cannot be made only with the Z_6 charges of the SM particles ($= 0$). If χ is much heavier than N, then the decay $\chi \to 3N^c + e^+ + e^-$, suppressed by $\tilde{e}_R$, E_R^c, and $\tilde{N}_R$ propagators, is possible.[1] Below, we consider the case of N being lighter than χ. The decay rate for $\chi \to 3N^c + e^+ + e^-$ is estimated as $\Gamma \sim g'^2|f|^4|h|^2 M_\chi^{11}/M_{\tilde{N}}^4 M_{\tilde{e}}^4 m_E^2$, but we require $M_\chi < 3m_N$ for successful two DM components in the universe. In this way, we have two DM components, χ and N. Now, we take the bino as the LNχ and the fermionic partner N as the LMP.

For this idea of two DM components to work, we must satisfy the following:

(i) The annihilation through Fig. 2 should be allowed. This will be one of the dominant sources of the positrons observed at PAMELA. Namely, $M_\chi + m_N > m_E + m_e$.

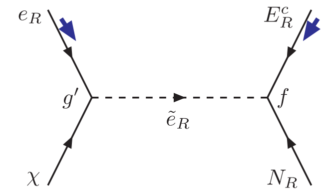

FIG. 2 (color online). A typical diagram for $N_R - \chi$ annihilation. $N_R N_R^c$ annihilation to e^+e^- is also possible.

(ii) χ or N should not decay by the diagram Fig. 2, and $M_\chi < m_N + m_E + m_e$ should be satisfied for $M_\chi > m_N$. The case $m_N > M_\chi$ turns out to be impossible.

(iii) The E decay is allowed by the interaction (2): $E_L \to e_R + \tilde{N}_R$, requiring $m_E > M_{\tilde{N}} + m_e$.

(iv) The bino decay $\chi \to 3N^c + e^+ + e^-$ is forbidden kinematically, $M_\chi < 3m_N + 2m_e$.

By the interaction (2), E_R^c by $\tilde{N}_R$ propagator ($E_R^c \to e_L^+ \tilde{N}_R^* \to e_L^+ N_R N_R$) and $\tilde{N}_R^*$ directly ($\tilde{N}_R^* \to 2N_R$ by (2)) decay to $2N_R + e_L^+$ and $2N_R$ for m_E, $M_{\tilde{N}} > 2m_N$, respectively. For the latter two body decay the total decay rate is

$$\Gamma(\tilde{N}_R^*) = \frac{|h|^2 M_{\tilde{N}} \beta_N}{16\pi}\left(1 - \frac{2m_N^2}{M_{\tilde{N}}^2}\right), \tag{7}$$

where $\beta_N = (1 - m_N^2/4M_{\tilde{N}}^2)^{1/2}$. The scalar partner of E, i.e. $\tilde{E}$, whose mass is definitely heavier than m_E, rapidly decays to eN. As a guide for the constraint on m_E, we use the direct search bound from LEP on the scalar lepton mass, $M_{\tilde{e}} \gtrsim 100$ GeV [19]. We take the LNχ mass of order 100 GeV, which would be a reasonable choice. Here, we assume that masses of N and $\tilde{N}_{R,L}$ are relatively small such that the needed kinematics are satisfied. Neglecting the electron mass, in Fig. 3 we plot the allowed region in the $M_\chi - m_E$ plane for a specific m_N. As far as the decay rate is large enough so that E and $\tilde{N}$ decayed before 1 s, these decays are not problematic in nucleosynthesis. Note that given an allowed phase space this is easily satisfied with not too small couplings f and h.

If χ and N each constitute 50% of the CDM density, the annihilation diagrams of Fig. 2 account for $\frac{1}{2}$ of possible

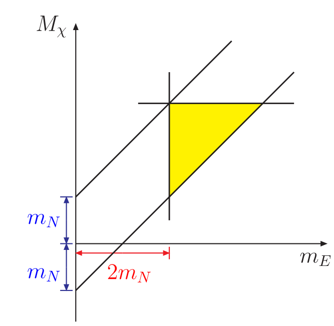

FIG. 3 (color online). In the $M_\chi - m_E$ plane, the kinematically allowed mass region is shaded for a typical mass value of m_N. For $M_\chi > 3m_N$, the decay $\chi \to 3N^c e^+ e^-$ is possible.

[1] Consider the Feynman diagram, $\chi \to Ee^+N^c$(by $\tilde{e}$ propagator) $\to (e^-\tilde{N})$(by E decay) $+ e^+N^c \to e^- + N^cN^c + e^+N^c$ (by $\tilde{N} \to N^cN^c$ of Eq. (2)).

encounters of χ and N: $\chi\chi$, NN, χN, and $N\chi$. Among these, the dominant contributions to the positron flux come from χN_R and $N_R N_R^c$ annihilation. In principle, the ratio of χ and N abundance is determined if high energy dynamics is completely known.[2]

Using the interaction (2), the cross section $d\sigma(\chi N \to e^+ E^-)/d\Omega$ of Fig. 2, in the center of momentum frame with the incident three momentum $\mathbf{p}$, is calculated in the small $|\mathbf{p}|$ and the large $M_{\tilde{E}}$ limit as

$$\frac{d\sigma}{d\Omega} \simeq \frac{|g'f|^2}{128\pi^2}\frac{1}{|\mathbf{p}|}\frac{m_\chi m_N \sqrt{s}\beta_E^4(1+m_E/\sqrt{s})^2}{[m_{\tilde{e}}^2 - m_\chi^2 + m_\chi\sqrt{s}\beta_E^2]^2} \times\left\{1+\left(\frac{2\sqrt{s}\beta_E^2}{m_{\tilde{e}}^2 - m_\chi^2 + m_\chi\sqrt{s}\beta_E^2} - \frac{1}{m_\chi} - \frac{\beta_E^2}{m_N(1+m_E/\sqrt{s})^2}\right)|\mathbf{p}|\cos\theta\right\}, \quad (8)$$

where θ is the angle between three momenta of χ and e^+, and $\beta_E^2 = 1 - m_E^2/s$ with $\sqrt{s} = M_\chi + m_N$. [We included the NN^c annihilation process also, which, however, is suppressed for a large $\tilde{E}_R$ mass.] A similar expression holds for the charge conjugated final states. Then, the velocity averaged cross section is calculated as

$$\langle\sigma v\rangle \simeq \frac{|g'f|^2}{32\pi^2}\frac{m_\chi m_N \beta_E^4(1+m_E/\sqrt{s})^2}{[m_{\tilde{e}}^2 - m_\chi^2 + m_\chi\sqrt{s}\beta_E^2]^2} + \mathcal{O}(v^2). \quad (9)$$

In Fig. 4, we present our estimate of the positron excess for typical values of $M_\chi = 200$ GeV, $m_N = 80$ GeV, $m_E = 200$ GeV, $M_{\tilde{E}} = 400$ GeV, and $M_{\tilde{e}} = 220$ GeV (thick green line), 250 GeV (blue dashed line), 280 GeV (brown dashed line). For a good fit, we need a small difference for $M_{\tilde{e}} - M_\chi$.

To compare with the observations, basically we use the astrophysical background flux given by $\Phi_{e^+}^{\rm bkg} = 4.5E^{0.7}/(1+650E^{2.3}+1500E^{4.2})$ and $\Phi_{e^-}^{\rm bkg} = 0.16E^{-1.1}/(1+11E^{0.9}+3.2E^{2.15}) + 0.70E^{0.7}/(1+110E^{1.5}+580E^{4.2})$ [20,21]. The deviation of the PAMELA data from this curve at low energy (< 10 GeV) can be explained by the solar modulation effect [5]. The calculation of the positron flux from a given particle physics model is well described in Refs. [22,23]. The positron flux is given by $\Phi_{e^+} = v_{e^+}\xi/4\pi$, where v_{e^+} is the velocity of the positron and ξ is the positron number density per unit energy, $\xi = dN_{e^+}/dE$. ξ is determined by the diffusion-loss equation using the various cosmic ray data as described in [22]. Under the steady state approximation, the solution of the diffusion-loss equation is given by a semiexact form

$$\Phi_{e^+} = \frac{Bv_{e^+}}{4\pi b(E)}\int_E^\infty dE' \sum_{i,j}\langle\sigma v\rangle_{i,j}\frac{\rho^2}{m_i m_j}\frac{dN}{dE'}I(\lambda_D(E,E')),$$

where $I(\lambda_D)$ is the halo function which has the halo model dependence but is independent from particle physics and $B \geq 1$ is a possible boost factor coming from the DM halo substructure [22].

So far, we have just assumed the contributions of χ and N to the total CDM density are comparable to each other. If the relic density of the dominant component among χ and N is 10 times larger than that of the other, the boost factor of about 100 would be needed. It could be explained by the DM clumpiness in the Galaxy.

In conclusion, it is likely that the LNχ of the MSSM cannot explain the high energy positron spectrum of the PAMELA/HEAT data [5,8], because the channel producing $e^\pm$ is relatively much suppressed. We proposed a SUSY model with the two dark matter particles χ and N, where additional charged lepton E and the $U(1)_R$ symmetry are also necessary for its completeness. It is the minimal extension of the MSSM, explaining the PAMELA/HEAT

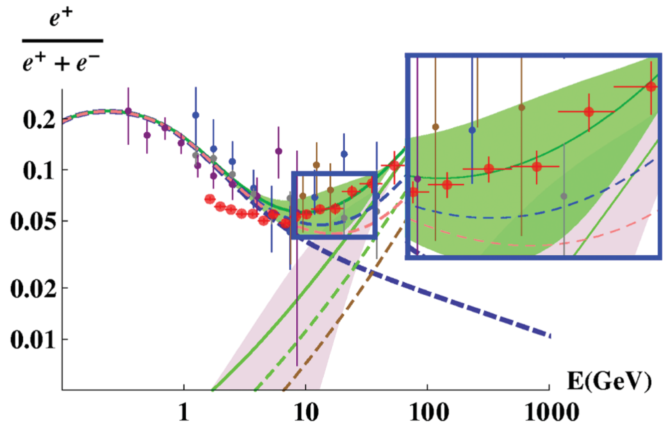

FIG. 4 (color). The positron fraction from our model with $M_\chi = 200$ GeV, $m_N = 80$ GeV, $m_E = 200$ GeV, $M_{\tilde{E}} = 400$ GeV, $M_{\tilde{e}} = 220$ GeV (thick green line), and $B = 7$. $M_{\tilde{e}}$ for 250 GeV (blue dashed line) and 280 GeV (brown dashed line) are also shown. The pink band is the primary positron fraction ($e^+_{\rm primary}/(e^+ + e^-)_{\rm total}$) coming from χN and NN annihilations and the green band is this positron excess on top of the astrophysical background (the thick dark-blue dashed line) [20,22]. The width of the band shows the uncertainty from the positron propagation model. The PAMELA data are the red dots [5], and the various small dots represent the observed positron cosmic ray data [8,24–26].

[2] In the case of the simple thermal production, the number density ratio $n_\chi/n_N \sim e^{-(M_\chi - m_N)/T_f}$ is expected, where T_f is the freeze-out temperature. For example, 10% mass difference could lead to the density ratio of ~ 0.1. However, nonthermal productions of DM might affect this result. For example, thermally produced axinos, which can explain the present relic density, subsequently decay into LNχ [4]. As a result, the number density of LNχ can be the same order of magnitude as that of N. Since the mass scale of N and LNχ is of weak scale, the resulting (co-)annihilation can lead to the needed amount of relic density for some parameter range. Moreover, since N is a SM singlet, it can easily couple to hidden sector fields, which can affect the thermal history of N. Because of such unknowns, we are satisfied just with showing that a concrete example with the assumption $\rho_N/\rho_\chi \sim 1$ can account for the positron excess.

data. It is shown that a wide range of the parameter space for two DM components is allowed in the model.

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