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Multicomponent dark matter in supersymmetric hidden sector extensions

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Most analyses of dark matter within supersymmetry assume the entire cold dark matter arising only from weakly interacting neutralinos. We study a new class of models consisting of $U(1)^n$ hidden sector extensions of the minimal supersymmetric standard model that includes several stable particles, both fermionic and bosonic, which can be interpreted as constituents of dark matter. In one such class of models, dark matter is made up of both a Majorana dark matter particle, i.e., a neutralino, and a Dirac fermion with the current relic density of dark matter as given by WMAP being composed of the relic density of the two species. These models can explain the PAMELA positron data and are consistent with the antiproton flux data, as well as the photon data from FERMI-LAT. Further, it is shown that such models can also simultaneously produce spin-independent cross sections which can be probed in CDMS-II, XENON-100, and other ongoing dark matter experiments. The implications of the models at the LHC and at the next linear collider (NLC) are also briefly discussed.

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I. INTRODUCTION

Recently several particle physics models have been constructed that connect the standard model (SM) to hidden sectors and lead to massive narrow vector boson resonances as well as other signatures which can be detected at colliders [1–3]. The connection to the hidden sector arises via mass mixings and kinetic mixings [1–6] and via higher dimensional operators. Models with the above forms of communication between the sectors also have important implications for dark matter [3,6,7] (for a review see [8,9]). In this work we show that multicomponent dark matter can arise from $U(1)^n$ extensions of the minimal supersymmetric standard model (MSSM) with Abelian hidden sectors which include hidden sector matter. Our motivation stems in part from the results of several dark matter experiments that have recently appeared. Thus the PAMELA Collaboration [10] has observed a positron excess improving previous results from HEAT and AMS experiments [11]. One possible explanation of such an excess is via the annihilation of dark matter in the galaxy [12]. Additionally, recent data from CDMS-II hints at the possibility of dark matter events above the background, and this will be explored further by the upgraded XENON experiment [13,14].

For a thermal relic, the PAMELA data and CDMS-II data taken together at face value do raise a theoretical puzzle if indeed both signals arise from the annihilation of cold dark matter. Thus most models which aim to explain the PAMELA positron excess do not give a significant number of dark matter events in the direct detection experiments currently operating. Conversely, models which can give a detectable signal in direct detection experiments typically do not explain the PAMELA data without the use of enormous so-called boost factors. As we will show here, this can be circumvented in models where the dark matter has several components. Thus, motivated in part by the recent cosmic anomalies we develop supersymmetric models which contain minimally a hidden Abelian sector broken at the sub-TeV scale where the mass generation of the hidden states involves nontrivial mixings with the field content of the electroweak sector of the minimal supersymmetric extension of the standard model leading to dark matter which can have several components which can be both bosonic and fermionic.

More specifically, in this work we go beyond the simple theoretical construction that thermal dark matter compatible with WMAP observations is composed of a single fundamental particle. There is no overriding principle that requires such a restriction, and nonbaryonic dark matter (DM) may indeed be constituted of several components, so in general one has $(\Omega h^2)_{\rm DM} = \sum_i (\Omega h^2)_{{\rm DM}i}$, where i refers to the various species of dark particles that can contribute to the total nonbaryonic $(\Omega h^2)_{\rm DM}$. In fact we already know that neutrinos do contribute to dark matter although their contribution is relatively small. Thus we propose here a new class of multicomponent cold dark matter models in Abelian $U(1)$ extensions of MSSM which can simultaneously provide an explanation of the PAMELA and WMAP data through a Breit-Wigner enhancement [12], while producing detectable signals for the direct searches for dark matter with CDMS/XENON and other dark matter experiments.

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A simultaneous satisfaction of the PAMELA positron excess and the satisfaction of WMAP relic density constraints can also occur if there is a nonthermal mechanism for the annihilation of dark matter with a wino lightest (R-parity odd) supersymmetric particle (LSP) [8,9,15–18]. However, a detectable spin-independent cross section in such a nonthermal framework does require that a pure wino is supplemented by a suitable admixture of Higgsino content as in the analysis of [19] and in [20], the later for a thermal relic. We remark that multiple $U(1)$ factors and its influence on dark matter have very recently been studied [20,21]. We also remark, some other works have recently looked at dark matter with more than 1 component [22]. The models proposed and analyzed here are very different from these.

The outline of the rest of the paper is as follows: In Sec. II we give a detailed description of the two models one of which is based on a $U(1)_X$ extension of the MSSM where $U(1)_X$ is a hidden sector gauge group with Dirac fermions in the hidden sector. This model allows for dark matter consisting of Dirac, Majorana, and spin zero particles. The second model is based on a $U(1)_X \times U(1)_C$ extension of MSSM , where $U(1)_C$ is a gauged leptophilic symmetry and $U(1)_X$, as before, is the hidden sector gauge group which also contains Dirac particles in the hidden sector. This model too has Dirac, Majorana, and spin zero particles as possible dark matter. In both cases we will primarily focus on the possibility that dark matter consists of Dirac and Majorana particles, and we will not discuss in detail the possibility of dark matter with bosonic degrees of freedom. In Sec. III we discuss the relic densities in the two component models. In Sec. IV we give an analysis of the positron, antiproton, and photon fluxes in the two models. In Sec. V we give an analysis of event rates for the proposed models for CDMS-II and for XENON-100. We give the analysis within the framework of supergravity grand unified models [23,24] defined by the parameters m_0, $m_{1/2}$, A_0, $\tan\beta$, and sign(μ) with nonuniversalities (NUSUGRA) defined by $\delta_{1,2,3}$ in the gaugino sector so that $U(1)_Y \times SU(2)_L \times SU(3)_C$ gaugino masses at the grand unified theory (GUT) scale are given by $\tilde{m}_i = m_{1/2}(1+\delta_i)$ $(i = 1, 2, 3)$ (see, e.g., [25] and references therein). We also discuss the possible new physics one might observe at the LHC (for a recent review see also [9]) and elsewhere for these models. Conclusions are given in Sec. VII.

II. MULTICOMPONENT HIDDEN SECTOR MODELS

A. Multicomponent $U(1)_X$ model

A $U(1)_X$ extension of the minimal supersymmetric standard model involves the coupling of a Stueckelberg chiral multiplet $S = (\rho + i\sigma, \chi_S, F_S)$ to vector supermultiplets X, B, where ρ is a real scalar and σ is an axionic pseudoscalar. Here X is the $U(1)_X$ vector multiplet which is neutral with respect to the SM gauge group with components $X = (X_\mu, \lambda_X, D_X)$, and B is the $U(1)_Y$ vector multiplet with components (B_μ, λ_B, D_B), where the components are written in the Wess-Zumino gauge. The chiral multiplet S transforms under both $U(1)_X$ and $U(1)_Y$ and acts as the connector sector between the visible and the hidden sectors. The total Lagrangian of the system is given by

$$\mathcal{L} = \mathcal{L}_{\rm MSSM} + \mathcal{L}_{U(1)_X} + \mathcal{L}_{\rm St}, \tag{1}$$

where $\mathcal{L}_{U(1)_X}$ is the kinetic energy piece for the X vector multiplet and $\mathcal{L}_{\rm St}$ is the supersymmetric Stueckelberg mixing between the X and the B vector multiplets so that [1,7] (see also [20,26,27])

$$\mathcal{L}_{\rm St} = \int d^2\theta d^2\bar{\theta}(M_1 X + M_2 B + S + \bar{S})^2, \tag{2}$$

where M_1 and M_2 are mass parameters. The Lagrangian of Eq. (1) is invariant under the $U(1)_Y$ and $U(1)_X$ gauge transformations, i.e., under

$$\begin{aligned} \delta_X X &= \zeta_X + \bar{\zeta}_X, \qquad & \delta_X S &= -M_1\zeta_X, \\ \delta_Y B &= \zeta_Y + \bar{\zeta}_Y, \qquad & \delta_Y S &= -M_2\zeta_Y, \end{aligned} \tag{3}$$

where ζ is an infinitesimal transformation chiral superfield. In component form we have for the Stueckelberg sector with $U(1)_X \times U(1)_Y$

$$\begin{aligned} \mathcal{L}_{\rm St} = &-\tfrac{1}{2}(M_1 X_\mu + M_2 B_\mu + \partial_\mu\sigma)^2 - \tfrac{1}{2}(\partial_\mu\rho)^2 \\ &- i\chi_S\sigma^\mu\partial_\mu\bar{\chi}_S + 2|F_S|^2 + \rho(M_1 D_X + M_2 D_B) \\ &+ \bar{\chi}_S(M_1\bar{\lambda}_X + M_2\bar{\lambda}_B) + \chi_S(M_1\lambda_X + M_2\lambda_B). \end{aligned} \tag{4}$$

In addition, one may include a supersymmetric kinetic mixing term between the $U(1)_X$ and $U(1)_Y$ gauge fields [7] leading to $\mathcal{L} = \mathcal{L}_{\rm MSSM} + \mathcal{L}_{U(1)_X} + \mathcal{L}_{\rm KM} + \mathcal{L}_{\rm St}$, where

$$\begin{aligned} \mathcal{L}_{U(1)_X} + \mathcal{L}_{\rm KM} = &-\frac{1}{4}X^{\mu\nu}X_{\mu\nu} - i\lambda_X\sigma^\mu\partial_\mu\bar{\lambda}_X + \frac{1}{2}D_X^2 \\ &- \frac{\delta}{2}X^{\mu\nu}B_{\mu\nu} - i\delta(\lambda_X\sigma^\mu\partial_\mu\bar{\lambda}_B \\ &+ \lambda_B\sigma^\mu\partial_\mu\bar{\lambda}_X) + \delta D_B D_X. \end{aligned} \tag{5}$$

One can also add additional D terms as in [7]. Both Stueckelberg and kinetic mixings of the gauge fields $U(1)_X$ and $U(1)_Y$ are constrained by the electroweak data [2]. As a consequence of the mixings, the extra gauge boson of the hidden sector couples with the standard model fermions and can become visible at colliders. The Lagrangian for matter interacting with the $U(1)$ gauge fields is given by

$$\begin{aligned} \mathcal{L}_{\rm matt} = \int d^2\theta d^2\bar{\theta}\sum_i[&\bar{\Phi}_i e^{2g_Y Q_Y B + 2g_X Q_X X}\Phi_i \\ &+ \bar{\Phi}_{{\rm hid},i} e^{2g_Y Q_Y B + 2g_X Q_X X}\Phi_{{\rm hid},i}], \end{aligned} \tag{6}$$

where the visible sector chiral superfields are denoted by

Φ_i (quarks, squarks, leptons, sleptons, Higgs, and Higgsinos of the MSSM) and the hidden sector chiral superfields are denoted by $\Phi_{\mathrm{hid},i}$. In the above, Q_Y is the hypercharge normalized so that $Q = T_3 + Q_Y$. As mentioned already, the SM matter fields do not carry any charge under the hidden gauge group and vice versa, i.e. $Q_X \Phi_i = 0$ and $Q_{\mathrm{SM}} \Phi_{\mathrm{hid}} = 0$. The minimal matter content of the hidden sector consists of a left chiral multiplet $\Phi_{\mathrm{hid}} = (\phi, f, F)$ and a charge conjugate $\Phi^c_{\mathrm{hid}} = (\phi', f', F')$ so that Φ_{hid} and Φ^c_{hid} have opposite $U(1)_X$ charges and form an anomaly-free combination. A mass M_ψ for the Dirac field ψ arises from an additional term in the superpotential $W_\psi = M_\psi \Phi \Phi^c$, where ψ is composed of f and f'. The scalar fields acquire soft masses of size m_0 from spontaneous breaking of supersymmetry by gravity mediation, and in addition acquire a mass from the term in the superpotential so that

$$m^2_\phi = m^2_0 + M^2_\psi = m^2_{\phi'}. \qquad (7)$$

After spontaneous breaking of the electroweak symmetry there would be mixing between the vector fields X_μ, B_μ, $A_{3\mu}$, where $A_{3\mu}$ is the third component of the $SU(2)_L$ field $A_{a\mu}$ $(a = 1, 2, 3)$. After diagonalization $V^T = (X, B, A_3)$ can be expressed in the terms of the mass eigenstates $E^T = (Z', Z, \gamma)$ as follows:

$$V_i = O_{ij} E_i, \qquad i, j = 1 - 3, \qquad E = (Z', Z, \gamma). \qquad (8)$$

The neutral vector mass squared matrix is of the form given in Ref. [1] of [6]. Further, the chiral fermions in the $S + \bar{S}$ multiplet together with the MSSM gauginos and Higgsinos will form a 6×6 neutralino mass matrix whose eigenstates are six neutralino states χ_a, $a = 1$–6, where we assume that the set $\chi^0_1 \ldots \chi^0_4$ is the regular set of neutralinos and χ^0_5, χ^0_6 are the two additional neutralinos that arise in the $U(1)_X$ extension. From the components λ_X, $\bar{\lambda}_X$ and χ_S, $\bar{\chi}_S$ that appear in Eq. (4), we can form two Majorana fields Λ_X and ψ_S as follows:

$$\Lambda_X = \begin{pmatrix} \lambda_{X\alpha} \\ \bar{\lambda}^{\dot{\alpha}}_X \end{pmatrix}, \qquad \psi_S = \begin{pmatrix} \chi_{\alpha,S} \\ \bar{\chi}^{\dot{\alpha}}_S \end{pmatrix}. \qquad (9)$$

These components combine with the MSSM gauginos and Higgsinos to form a 6×6 neutralino mass matrix whose eigenstates are the six neutralinos χ_a $(a = 1$–$6)$. Thus Λ_X and ψ_S can be expanded as linear combination of χ_a, i.e.,

$$\begin{aligned} \Lambda_X &= R_{1a} \chi_a, \qquad a = 1 - 6, \\ \psi_S &= R_{2a} \chi_a, \qquad a = 1 - 6, \end{aligned} \qquad (10)$$

where R is the unitary matrix which diagonalizes the 6×6 neutralino mass matrix. Further the CP even Higgs sector is extended by the additional state ρ [1]. The results outlined here give the following types of interactions:

(i) There are interactions of the Dirac fermion in the hidden sector with the standard model particles via Z, Z', γ interactions. Thus, the Dirac dark matter can annihilate into standard model particles via exchange of Z, Z', γ in the early universe and in the galaxy. Depending on which of the two, Dirac or Majorana, is the heavier one may have Dirac particles annihilating into Majoranas or the Majorana particles annihilating into Dirac fermions in the galaxy:

$$\bar{\psi}\psi \to \chi\chi \quad \text{or} \quad \chi\chi \to \bar{\psi}\psi. \qquad (11)$$

(ii) In addition to the above we have fermion-neutralino-sfermion couplings in the hidden sector as given by Eq. (6). Thus interactions of the type $\bar{\psi}\chi_a\phi + \mathrm{H.c.}$, etc. can produce decays such as $\phi \to \psi + \chi_a$ if they are kinematically allowed.

(iii) The scalar field ρ is CP even and mixes with the MSSM Higgs fields. Through these mixings ρ has couplings to the SM fermions and through these couplings it can decay into the SM fermions.

It is instructive to list all the new particles in this $U(1)_X$ model as summarized below:

$$\begin{aligned} &\text{New particles of the } U(1)_X \text{ model} \\ &\text{spin 0: } \rho, \phi, \phi', \\ &\text{spin } \tfrac{1}{2}\text{: } \psi, \chi^0_5, \chi^0_6, \\ &\text{spin 1: } Z'. \end{aligned} \qquad (12)$$

We assume that the lightest R-parity odd particle (LSP) is the least massive neutralino ($\chi^0 = \chi^0_1 \equiv \chi$) and resides in the visible sector and thus the masses of χ^0_5, χ^0_6 are larger than the LSP χ^0 mass, and consequently χ^0_5, χ^0_6 are unstable and decay into SM particles and χ^0. The bosons Z' and ρ are unstable and decay into SM fermion pairs $f\bar{f}$ with the decay of the ρ going dominantly through the process $\rho \to b\bar{b}$ or $\rho \to t\bar{t}$ if $m_\rho > 2m_t$. The remaining three particles ψ, ϕ, ϕ' are all milli charged and, consequently, at least one of them is stable. If we assume m_ϕ, $m_{\phi'} > M_\psi$, at least ψ is always stable and the other two may or may not be stable. These along with the LSP give rise to various possible candidates for dark matter. Thus, depending on the relative masses of the Majorana, Dirac, and spin 0 particles there are three possibilities for the constituents of dark matter as outlined below.

Two component dark matter: Majorana + Dirac.—This model arises as follows: consider the case where $m_\phi > M_\psi + M_\chi$. In this case the decays $\phi, \phi' \to \psi + \chi^0$, will occur and ϕ, ϕ' will be unstable. Thus ψ is stable and so is χ under the assumption of R parity conservation. Consequently, we will have two dark matter particles; namely, one a Majorana which is the LSP in the visible sector and the other a Dirac in the hidden sector. The Majorana and Dirac particles once created will annihilate as follows:

$$\psi + \bar{\psi} \to Z, Z', \gamma \to \mathrm{SM} + \mathrm{SM}', \qquad (13)$$

$$\chi + \chi \rightarrow (s\colon Z', Z, h, H, A, \rho), \quad (t/u\colon \tilde{f}_a, \chi_i, \chi_k^{\pm}) \rightarrow \mathrm{SM} + \mathrm{SM}', \tag{14}$$

where s : and t/u : refer to s and t or u channel exchanges. In addition to Eq. (14) there are coannihilation processes which contribute to the relic density. Since both ψ and χ are stable, the total relic density of dark matter will be the sum of the relic densities for the two, the sum being constrained by the WMAP data. These constraints are discussed further in Sec. III.

Three component dark matter: Dirac and two spin 0 particles.—Suppose the mass of χ is larger than the sum of the masses of the Dirac plus the scalar ϕ, i.e., $M_\chi > M_\psi + m_\phi$. In this case the decay $\chi \rightarrow \phi + \psi, \phi' + \psi$ will occur and, consequently, χ is unstable. On the other hand, ϕ, ϕ' and ψ are stable since they cannot decay into anything else. Thus, here we have three dark matter particles: one Dirac, and the other two spin 0. Processes that lead to the annihilations of these particles are those in Eq. (13) for ψ, and also for ϕ and ϕ', they are similar to those in Eq. (13), i.e., $\phi + \phi^*$, $\phi' + \phi'^* \rightarrow \gamma, Z, Z' \rightarrow \mathrm{SM} + \mathrm{SM}'$. In this three component dark matter model all the components reside in the hidden sector and thus their couplings to the standard model particles are extra weak. Consequently, they will have very small spin-independent cross sections in direct detection experiments. For this reason, this class of models is less preferred compared to the two component model.

Four component dark matter: Majorana, Dirac, and two spin 0 particles.—Finally, we consider the case when either of the following two situations occur: (i) $M_\chi > M_\psi$, $m_\phi < M_\chi < M_\psi + m_\phi$, (ii) $M_\chi < m_\phi < M_\chi + M_\psi$. In these cases all four particles, one Majorana, one Dirac, and two spin 0 particles, are stable and thus are possible dark matter candidates. These particles will annihilate to the SM particles as in Eqs. (13) and (14) and for ϕ and ϕ' via processes in the three component dark matter model as described above. This model is in many ways similar to the two component model and like the two component model this model too should lead to detectable signals in experiments for the direct detection of dark matter.

B. Multicomponent leptophilic $U(1)_X \times U(1)_C$ model

We discuss now another model which contains two additional Abelian vector bosons where one of the extra bosons is leptophilic. Leptophilic Z's have a long history [28] and have been revisited [29] over the recent past in the context of dark matter. Here we will consider a $U(1)_X \times U(1)_C$ model where the $U(1)_X$ as before is in the hidden sector, and $U(1)_C$ is a leptophilic symmetry. As in the $U(1)_X$ model, we also assume that the hidden sector has a pair of Dirac fermions ψ and $\bar{\psi}$ which are charged under $U(1)_X$ but are neutral under the standard model gauge group and under $U(1)_C$. Regarding $U(1)_C$ we assume it to be $L_e - L_\mu$, i.e., a difference of family-lepton numbers, which is anomaly free, and can be gauged. The corresponding gauge field C_μ couples only to e, μ families and nothing else. The total Lagrangian in this case is

$$\mathcal{L} = \mathcal{L}_{\mathrm{MSSM}} + \mathcal{L}_{U(1)^2} + \mathcal{L}_{\mathrm{St}}, \tag{15}$$

where $\mathcal{L}_{U(1)^2}$ is the kinetic energy for the X and C multiplets and for $\mathcal{L}_{\mathrm{St}}$ we assume the following form:

$$\mathcal{L}_{\mathrm{St}} = \int d^2\theta d^2\bar{\theta}(M_1 C + M_2' X + M_3' B + S + \bar{S})^2 + \int d^2\theta d^2\bar{\theta}(M_1' C + M_2 X + M_3'' B + S' + \bar{S}')^2, \tag{16}$$

where C is the $U(1)_{L_e - L_\mu}$ vector multiplet with components (C_μ, λ_C, D_C) and X and B are the $U(1)_X$ and $U(1)_Y$ multiplets as discussed before. The gauge transformations under $U(1)_C$, $U(1)_X$, and $U(1)_Y$ are

$$\begin{aligned} &\delta_C C = \zeta_C + \bar{\zeta}_C, && \delta_C S = -M_1 \zeta_C, \\ &\delta_C S' = -M_1' \zeta_C && \delta_X X = \zeta_X + \bar{\zeta}_X, \\ &\delta_X S = -M_2' \zeta_X, && \delta_X S' = -M_2 \zeta_X, \\ &\delta_Y B = \zeta_Y + \bar{\zeta}_Y, && \delta_Y S = -M_3' \zeta_Y, \\ &\delta_Y S' = -M_3'' \zeta_Y, && \end{aligned} \tag{17}$$

where ζ_C, ζ_X, ζ_Y, etc. are the infinitesimal transformation chiral superfields. The quantities M_1, M_2, M_1', M_2', M_3', and M_3'' are the mass parameters. In the vector boson sector $\mathcal{L}_{\mathrm{St}}$ assumes the form

$$\mathcal{L}_{\mathrm{St}} = -\tfrac{1}{2}(M_1 C_\mu + M_2' X_\mu + M_3' B_\mu + \partial_\mu \sigma)^2 - \tfrac{1}{2}(M_1' C_\mu + M_2 X_\mu + M_3'' B_\mu + \partial_\mu \sigma')^2. \tag{18}$$

The mass^2 matrix in the vector boson sector in the basis $(C^\mu, X^\mu, B^\mu, A^{3\mu})$ is given by

$$\begin{pmatrix} M_1^2 + M_1'^2 & M_1 M_2' + M_1' M_2 & M_1 M_3' + M_1' M_3'' & 0 \\ M_1 M_2' + M_1' M_2 & M_2^2 + M_2'^2 & M_2' M_3' + M_2 M_3'' & 0 \\ M_1 M_3' + M_1' M_3'' & M_2' M_3' + M_2 M_3'' & M_3'^2 + M_3''^2 + M_Y^2 & -M_Y M_W \\ 0 & 0 & -M_Y M_W & M_W^2 \end{pmatrix}, \tag{19}$$

where $M_W = g_2 \cdot v/2$ is the W boson mass and $M_Y = M_W \tan\theta_W = g_Y \cdot v/2$, and where θ_W is the weak angle. The dynamics of the model of Eq. (19) is rather involved. We will focus, therefore, on a simpler version of this more general case where we neglect the mixings with B_μ, i.e., we set $M_3' = M_3'' = 0$. Inclusion of these coupling in the analysis would not drastically change the analysis or the conclusions of this work as long as we keep the mixing parameters $M_3'/M_{1,2}$, $M_3''/M_{1,2}$ very small. After neglecting the mixings with B_μ, the mass^2 matrix is block diagonal and so we can diagonalize the top left hand corner 2×2 mass matrix independent of the standard model sector. We are interested in the limit of small mixing between $U(1)_X$ and $U(1)_C$ and thus consider[1]

$$M_1', M_2' \ll M_1, M_2. \tag{20}$$

In the above approximation the eigenvalues of this mass matrix are

$$\begin{aligned} M_{Z'}^2 &\simeq M_2^2 + M_2'^2 - \Delta_{M^2},\\ M_{Z''}^2 &\simeq M_1^2 + M_1'^2 + \Delta_{M^2},\\ \Delta_{M^2} &\simeq \frac{(M_1M_2' + M_1'M_2)^2}{(M_1^2 + M_1'^2 - M_2^2 - M_2'^2)}. \end{aligned} \tag{21}$$

The corresponding mass eigenstates are Z' and Z'', where

$$\begin{aligned} C_\mu &= \cos\theta_X Z_\mu'' - \sin\theta_X Z_\mu',\\ X_\mu &= \sin\theta_X Z_\mu'' + \cos\theta_X Z_\mu',\\ \tan\theta_X &\simeq \frac{M_1M_2' + M_1'M_2}{M_1^2 + M_1'^2 - M_2^2 - M_2'^2}. \end{aligned} \tag{22}$$

Because of Eq. (20) $\tan\theta_X \ll 1$. In the above, the Dirac fermions in the hidden sector have no couplings with the photon and are electrically neutral. However, by a small mixing of X_μ with B_μ in Eq. (18), we can generate a milli charge for the Dirac particles in the hidden sector consistent with all electroweak data.

We discuss now the gaugino/chiral fermions in the extra $U(1)$ sectors which arise from the superfields C, X, $S+\bar{S}$, $S'+\bar{S}'$. From the gaugino components λ_C, $\bar{\lambda}_C$, λ_X, $\bar{\lambda}_X$, and from the chiral fermion components in the extra $U(1)$ sectors χ_S, $\bar{\chi}_S$, $\chi_{S'}$, $\bar{\chi}_{S'}$, one can construct four component Majorana spinors two of which are exhibited in Eq. (9) and the remaining two are given by

$$\Lambda_C = \begin{pmatrix} \lambda_{C\alpha} \\ \bar{\lambda}_C^{\dot\alpha} \end{pmatrix}, \qquad \psi_{S'} = \begin{pmatrix} \chi_{\alpha,S'} \\ \bar{\chi}_{S'}^{\dot\alpha} \end{pmatrix}. \tag{23}$$

The neutralino mass matrix in the $[U(1)_X \times U(1)_C] \times [SU(3)_C \times SU(2)_L \times U(1)_Y]$ model takes a block diagonal form

$$\begin{bmatrix} U(1)_X \times U(1)_C \text{ sector} & 0_{4\times4} \\ 0_{4\times4} & \text{MSSM sector} \end{bmatrix}_{8\times8}. \tag{24}$$

Thus, the Stueckelberg mass generation produces a mass matrix in the hidden gaugino/chiral fermion sector which is decoupled from the neutralino mass matrix in the visible sector. Specifically in the 4 component notation the gaugino/chiral fermion mass matrix in the $U(1)_X \times U(1)_C$ sector is given by

$$L^{\text{mass}}_{U(1)_X\times U(1)_C} = -\begin{pmatrix} \bar\psi_S \\ \bar\psi_{S'} \\ \bar\Lambda_C \\ \bar\Lambda_X \end{pmatrix}^T \begin{pmatrix} 0 & 0 & M_1 & M_2' \\ 0 & 0 & M_1' & M_2 \\ M_1 & M_1' & 0 & 0 \\ M_2' & M_2 & 0 & 0 \end{pmatrix} \begin{pmatrix} \psi_S \\ \psi_{S'} \\ \Lambda_C \\ \Lambda_X \end{pmatrix}. \tag{25}$$

In the diagonalized basis we can label the extra neutralinos by χ_5^0, χ_6^0, χ_7^0, χ_8^0. Since the hidden sector and the neutralinos of the visible sector are decoupled, the diagonalization of the neutralinos in the visible sector, i.e., of χ_i^0, $(i = 1-4)$ is not affected. Further, as for the case of the $U(1)_X$ model, it is instructive to list all the new particles in this $U(1)_X \times U(1)_C$ model as summarized below:

$$\begin{aligned} &\text{New particles of } U(1)_C \times U(1)_X \text{ model}\\ &\text{spin 0: } \rho, \rho', \phi, \phi',\\ &\text{spin } \tfrac{1}{2}\text{: } \psi, \chi_5^0, \chi_6^0, \chi_7^0, \chi_8^0,\\ &\text{spin 1: } Z', Z''. \end{aligned} \tag{26}$$

We discuss now the stability of the new particles in this model. As before we assume that the mass of ϕ (and of ϕ') is larger than the mass of ψ. Thus ψ will be stable since it cannot decay into anything. If kinematically allowed the fields ϕ and ϕ' can decay only via the process ϕ, $\phi' \to \psi + \chi^0$ as in the $U(1)_X$ model. Of the remaining fields obviously Z' and Z'' are unstable as they decay into $e\bar{e}$, $\mu\bar{\mu}$, $\nu_e\bar{\nu}_e$, $\nu_\mu\bar{\nu}_\mu$ as well as into $\psi\bar{\psi}$ depending on the mass of ψ. As already noted, a small milli charge can develop for the hidden sector matter via small couplings of the B_μ and X_μ fields. The phenomenology of such models will be very similar to the one we are discussing here.

The extra neutralinos of Eq. (26) can also be all unstable. Thus, Λ_C couples with leptons-sleptons (e, $\tilde{e}$, etc.) via coupling of the type $\bar\Lambda_C e_L \tilde{e}_L^*$, etc. and after diagonalization of the gaugino/chiral fermion mass matrix all the χ_k^0 ($k =$ 5–8) will have coupling with leptons-sleptons of the type indicated. Further, two of the χ_k^0 have roughly a mass of size M_1 while the remaining two have roughly a mass of size M_2. Thus, if $M_1, M_2 > m_{\chi^0}$, which is what is assumed in this work, all the neutralinos of the hidden sector will be unstable and decay into final states of the type $e\bar{e}\chi^0$, $\mu\bar{\mu}\chi^0$, etc. Regarding the field ρ, there is an interaction of type

$$M_1 g_C \rho(\tilde{f}^* Q_C^f \tilde{f}), \qquad f = e, \mu. \tag{27}$$

With this interaction ρ will decay as follows: $\rho \to \tilde{f}^*\tilde{f} \to f\bar{f}\chi^0\chi^0 (f = e, \mu)$ provided this process is kinematically allowed which we assume is the case. A similar situation occurs for the case of ρ'. Additionally, if there is a mixing with B_μ in the Stueckelberg sector then, as in the analysis of the $U(1)_X$ model, the fields ρ and ρ' will mix

[1] Note these mass terms M_1, M_2 are different than those considered in Sec. II A.

with the Higgs sector and can have decays of the type $\rho \to b\bar{b}$, $\rho' \to b\bar{b}$, etc. Thus, in the end we are left with a similar set of possibilities for dark matter as in the $U(1)_X$ model, i.e., (i) a two component model with ψ and χ^0, (ii) a three component model with ψ, ϕ, ϕ', and (iii) a four component model with ψ, ϕ, ϕ', and χ^0. However, as in the $U(1)_X$ case we will focus on the two component model consisting of Dirac and Majorana dark particles.

We assume $M^2_{Z''} \gg M^2_{Z'}$ and that the annihilation of dark matter occurs close to the Z' pole for reasons that will become apparent shortly. As a consequence, the annihilation of dark matter in the early universe and in the galaxy is controlled by the Z' pole and the effect of the Z'' pole on the analysis is essentially negligible. The basic interaction of C_μ and of X_μ with matter is given by

$$\mathcal{L}_{\rm int} = g_X Q_X \bar{\psi}\gamma^\mu \psi X_\mu + g_C Q_C^f \bar{f}\gamma^\mu f C_\mu, \qquad (28)$$

where f runs over e and μ families and where $Q_C^e = -Q_C^\mu$. In the mass diagonal basis the interaction of Eq. (28) assumes the form

$$\begin{aligned}\mathcal{L}_{\rm int} = {}& (g_X Q_X \bar{\psi}\gamma^\mu \psi \cos\theta_X - g_C Q_C^f \bar{f}\gamma^\mu f \sin\theta_X) Z'_\mu \\ & + (g_X Q_X \bar{\psi}\gamma^\mu \psi \sin\theta_X + g_C Q_C^f \bar{f}\gamma^\mu f \cos\theta_X) Z''_\mu.\end{aligned} \qquad (29)$$

The interaction of Eq. (29) leads to the annihilation of $\psi\bar{\psi}$ into e^+e^- and $\mu^+\mu^-$ via the Z', Z'' poles for which we assume a Breit-Wigner form. Thus, the $\psi\bar{\psi} \to f\bar{f}$ annihilation cross section takes the form

$$\begin{aligned}\sigma_{\psi\bar{\psi}\to f\bar{f}} = {}& a_\psi |(s - M^2_{Z'} + i\Gamma_{Z'} M_{Z'})^{-1} \\ & - (s - M^2_{Z''} + i\Gamma_{Z''} M_{Z''})^{-1}|^2,\end{aligned} \qquad (30)$$

$$\begin{aligned}a_\psi = {}& \frac{\beta_f (g_X g_C Q_X Q_C^f \sin(2\theta_X))^2}{64\pi s \beta_\psi}\Big[s^2\Big(1 + \frac{1}{3}\beta_f^2\beta_\psi^2\Big) \\ & + 4M_\psi^2(s - 2m_f^2) + 4m_f^2(s + 2M_\psi^2)\Big],\end{aligned} \qquad (31)$$

where $\beta_{f,\psi} = (1 - 4m^2_{f,\psi}/s)^{1/2}$. The relevant partial Z' decay widths are given by

$$\Gamma(Z' \to f\bar{f}) = (g_C Q_C^f \sin\theta_X)^2 \frac{M_{Z'}}{12\pi}, \qquad f = e, \mu, \qquad (32)$$

$$\begin{aligned}\Gamma(Z' \to \psi\bar{\psi}) = {}& (g_X Q_X \cos\theta_X)^2 \frac{M_{Z'}}{12\pi}\Big(1 + \frac{2M_\psi^2}{M^2_{Z'}}\Big) \\ & \times \Big(1 - \frac{4M_\psi^2}{M^2_{Z'}}\Big)^{1/2} \Theta(M_{Z'} - 2M_\psi),\end{aligned} \qquad (33)$$

and similarly for the partial decay widths of the Z'' with $M_{Z'} \to M_{Z''}$ and $-\sin\theta_X \to \cos\theta_X$ in Eq. (32) and $\cos\theta_X \to \sin\theta_X$ in Eq. (33).

A constraint on g_C comes from the contribution of the Z' and Z'' to $g_\mu - 2$ [30]. Their exchange gives

$$\Delta(g_\mu - 2) = \frac{g_C^2 m_\mu^2}{24\pi^2}\Big[\frac{\sin^2\theta_X}{M^2_{Z'}} + \frac{\cos^2\theta_X}{M^2_{Z''}}\Big]. \qquad (34)$$

Using the current error [30] of $\Delta(g_\mu - 2) = 1.2 \times 10^{-9}$ in the determination of $g_\mu - 2$ and assuming θ_X is small, one finds the following constraint on α_C:

$$\alpha_C \lesssim 0.001\Big(\frac{M_{Z''}}{300\ \text{GeV}}\Big)^2, \qquad (35)$$

where $\alpha_C = g_C^2/4\pi$. We note that if the mixing angle θ_X is small, the decay width of $Z' \to f\bar{f}$ ($f = e, \mu$) and of $Z'' \to \psi\bar{\psi}$ will be narrow while the decay width of $Z'' \to f\bar{f}$ ($f = e,\ \mu$) and of $Z' \to \psi\bar{\psi}$ will be of normal size. However, when $M_\psi \simeq M_{Z'}/2$ the Z' decay width into $\psi\bar{\psi}$ will also be small due to the kinematic suppression factor $\{1 - [(4M_\psi^2)/M^2_{Z'}]\}^{1/2}$. In this case we will have the total width of the Z' to be rather narrow. Thus for annihilation near the Breit-Wigner pole we will have a large enhancement of $\langle\sigma v\rangle$ due to the narrowness of the Z' [12]. It was shown in the analysis of Feldman-Liu-Nath in [12] that near the Breit-Wigner pole such annihilations allow one to fit the relic density as well as allow an enhancement of $\langle\sigma v\rangle$ in the galaxy. We note that while Z' decay width is very small this is not necessarily the case for Z'' which can decay into $e\bar{e}$, $\mu\bar{\mu}$, $\nu_e\bar{\nu}_e$, $\nu_\mu\bar{\nu}_\mu$ with normal strength. Thus, neglecting the contribution of $Z'' \to \psi\bar{\psi}$ which is small due to the $\sin^2\theta_X \sim \epsilon^2$ suppression, one finds the total width of Z'' to be $\Gamma_{Z''} \simeq \cos^2\theta_X \alpha_C M_{Z''}$. We will see in Sec. (4) that the α_C needed in the analysis of the relic density is relatively small compared to normal electroweak coupling and, consequently, the width of Z'' though significantly larger than the Z' width is still relatively small compared to what one might expect for a Z' in a GUT model and certainly much smaller than the width for a Z' arising as a Kaluza-Klein excitation in the compactification of an extra dimension [31,32]. Finally, the annihilation of the Dirac particles in the early universe goes by the processes

$$\psi\bar{\psi} \to Z', Z'' \to e^+e^-, \mu^+\mu^-, \nu_e\bar{\nu}_e, \nu_\mu\bar{\nu}_\mu, \qquad (36)$$

which is to be contrasted with the processes Eq. (13) in the $U(1)_X$ model.

III. RELIC DENSITY IN A TWO COMPONENT MODEL

Here we discuss the relic density in models with two components. A general analysis requires solving the Boltzmann equations in a Friedmann-Robertson-Walker universe [33,34], and includes coannihilations [35] and an accurate integration over pole regions. As in the MSSM alone, one will generally encounter the Z and Higgs poles [36] and these need to be treated with care. The number changing processes include

$$\psi\bar{\psi} \leftrightarrow \mathrm{SM}\,\mathrm{SM}', \qquad \psi\bar{\psi} \leftrightarrow \chi\chi, \qquad \chi\chi \leftrightarrow \mathrm{SM}\,\mathrm{SM}'. \tag{37}$$

Note that the process $\bar{\psi}\chi \leftrightarrow \mathrm{SM}\,\mathrm{SM}'$ is not allowed since $\bar{\psi}\chi$ connect only to ϕ and ϕ', neither of which can connect to the standard model particles. For the simplest two component model with dark matter particles ψ, χ, with the assumption that $M_\psi > M_\chi$ the only relevant processes in the annihilation of $\psi\bar{\psi}$ are $\psi\bar{\psi} \to f\bar{f}$, $\chi\chi$ final states. Since ψ is heavier than χ its freeze-out occurs earlier (at a higher T) than for χ. Thus, the Boltzmann equations for n_ψ (which includes fermions and antifermions) and for n_χ for the $U(1)_X$ and for the $U(1)_X \times U(1)_C$ two component models are given by

$$\frac{dn_\psi}{dt} = -3Hn_\psi - \frac{1}{2}\langle\sigma v\rangle_{\psi\bar{\psi}}(n_\psi^2 - n_{\psi,\mathrm{eq}}^2), \tag{38}$$

$$\frac{dn_\chi}{dt} = -3Hn_\chi - \langle\sigma v\rangle_{\chi\chi}(n_\chi^2 - n_{\chi,\mathrm{eq}}^2) + \frac{1}{2}\langle\sigma v\rangle_{\psi\bar{\psi}\to\chi\chi}(n_\psi^2 - n_{\psi,\mathrm{eq}}^2). \tag{39}$$

Here $\langle\sigma v\rangle_{\psi\bar{\psi}}$ refers to $\psi\bar{\psi} \to f\bar{f}$, $\chi\chi$, and $\langle\sigma v\rangle_{\chi\chi}$ stands for $\langle\sigma v\rangle_{\chi\chi\to\mathrm{SM}\,\mathrm{SM}'}$. For the spin averaged cross section for the Dirac case, the extra factor of $1/2$ is to account for the fact that we are dealing with a Dirac fermion. The number densities are n_ψ, n_χ and $n_{\psi,\mathrm{eq}}, n_{\chi,\mathrm{eq}}$ are their values at equilibrium, i.e., $n_{(\psi,\chi),\mathrm{eq}} \simeq g_{(\psi,\chi)}(M_{(\psi,\chi)}T)/2\pi)^{3/2} \times \exp(-\frac{M_{(\psi,\chi)}}{T})$, where $g_\psi = 4$ and $g_\chi = 2$. Since the two dark matter particles are sub-TeV in mass, they will freeze-out at temperatures that are not drastically different. One can solve the Boltzmann equation for ψ with the appropriate boundary conditions to compute the freeze-out temperature T_f^ψ and the relic density of ψ at the current temperatures. To compute the freeze-out temperature T_f^χ for the particles χ, one uses solutions for n_ψ as computed from the Boltzmann equation for ψ as input in the Boltzmann equation for χ keeping in mind that $n_{\psi,\mathrm{eq}}$ in the χ Boltzmann equation can be neglected since we are below the freeze-out temperature for ψ. It is difficult to get a closed form solution of Eq. (39) for n_χ and thus in general the analysis must be done numerically for $\Omega_\chi h^2$. However, it turns out that for both the $U(1)_X$ and the $U(1)_X \times U(1)_C$ models the contribution of the term proportional to n_ψ^2 in Eq. (39) is rather suppressed and it is a good approximation to neglect this term for both models. In this case, one has

$$(\Omega h^2)_{\mathrm{WMAP}} = (\Omega_\psi h^2)_0 + (\Omega_\chi h^2)_0 \simeq \frac{C_\psi}{J_0^\psi} + \frac{C_\chi}{J_0^\chi}, \tag{40}$$

where

$$C_\chi \simeq \frac{1.07 \times 10^9\ \mathrm{GeV}^{-1}}{\sqrt{g^*(\chi)}M_{\mathrm{pl}}}, \qquad C_\psi \simeq 2 \times \frac{1.07 \times 10^9\ \mathrm{GeV}^{-1}}{\sqrt{g^*(\psi)}M_{\mathrm{pl}}}, \tag{41}$$

$$J_0^\chi = \int_0^{x_f^\chi}\langle\sigma v\rangle_{\chi\chi}dx, \qquad J_0^\psi = \int_0^{x_f^\psi}\langle\sigma v\rangle_{\psi\bar{\psi}}dx, \tag{42}$$

and where $g^*(\psi, \chi)$ denotes the effective degrees of freedom at the freeze-out of ψ, χ, respectively. The analysis leading to Eqs. (38) and (39) is easily extended to include coannihilations. The analysis can easily be reversed if the Majorana is heavier than the Dirac. Denoting $\rho_{\odot,\psi}$, $\rho_{\odot,\chi}$ as the local density of each dark matter kind in the halo, one can assume

$$\rho_{\odot,\psi}/\rho_{\odot,\chi} \sim (\Omega_\psi h^2)_0/(\Omega_\chi h^2)_0. \tag{43}$$

However, the ratios need not be the same. The local halo densities are also constrained such that $\rho_{\odot,\psi} + \rho_{\odot,\chi} = \rho_{\odot,\mathrm{total}} \simeq (0.35\text{–}0.45)\ \mathrm{GeV\,cm^{-3}}$. For the calculation near the Z' pole, we use the analysis of [6] which follows the techniques of [36]. Indeed the analytic techniques developed in [3,36] have been cross-checked with independent codes. For the $U(1)_X$ model, the decay branching ratios are substantially less hadronic and more leptonic than for the annihilations via the Z boson exchange [37]. For the $U(1)_X \times U(1)_C$ model, the decays of the Z', Z'' are purely leptonic. These leptophilic decay patterns for the extra Z's help to explain the PAMELA positron excess without recourse to large *ad hoc* boost factors.

IV. POSITRON, ANTIPROTON, AND PHOTON FLUXES IN THE $U(1)_X$ AND $U(1)_X \times U(1)_C$ MODELS

An excess of positrons, antiprotons, and photons over the cosmic background is a possible indicator for annihilating dark matter in the galaxy as was pointed out early on [38]. In our multicomponent supersymmetric models contributions from the Majorana component to the fluxes are negligible (suppressed by an order of magnitude or more) and essentially the entire effect arises from the Dirac component. The positron flux Φ_{e^+} arising from the annihilation of dark matter (DM) particles is [39,40]

$$\Phi_{e^+} = \frac{\eta v_{e^+}}{4\pi b(E)}\frac{\rho_\odot^2}{M_{\mathrm{DM}}^2}\int_E^{M_{\mathrm{DM}}}\sum_f\langle\sigma v\rangle_{f,\mathrm{halo}}\frac{dN_{e^+}^f}{dE'}B_{\tilde{e}}I_{(E,E')}dE'. \tag{44}$$

Here M_{DM} is the mass of the dark matter particle. In the above $\eta = (1/2, 1/4)$ for (Majorana, Dirac) cases, respectively, $B_{\tilde{e}}$ is a boost factor which may arise as a consequence of dark matter substructure, or a local clump. Recent N-body simulations show that it is unlikely that large dark matter clumps exist within the halo of our galaxy [41] and thus the use of large clump factors in flux analyses appear unreasonable [42].

The other parameters that enter Eq. (44) are as follows: $I_{(E,E')}$ is the halo function and we parametrize it in some of the standard forms adopted in the literature with the appropriate diffusion models [39] using the standard profiles [43]. The positron velocity is $v_{e^+} \sim c$, and the energy loss function $b(E)$ has the form $b(E) = E_0(E/E_0)^2/\tau_E$, where $\tau_E \sim (1\text{–}2)10^{16}$ [s], with E in [GeV] and $E_0 \equiv 1$ GeV. We use the GALPROP background estimate of [44] fit in [45]; modifications of the background estimates require either smaller or larger mass splitting at the pole which can range from the order of a GeV to the order of tens of MeV, depending on the assumed astrophysical background and the level of clumpiness of the signal. Further, in Eq. (44) $\langle\sigma v\rangle_{\text{halo}}$ is the velocity averaged cross section in the *halo of the galaxy*. We note that $\langle\sigma v\rangle_{\text{halo}}$ may be significantly different than $\langle\sigma v\rangle_{x_f}$ at the epoch of freeze-out. Thus, as emphasized in [12] a replacement of $\langle\sigma v\rangle_{\text{halo}}$ by $\langle\sigma v\rangle_{\text{freeze-out}}$, as is often done, is inaccurate and can lead to significant errors in the positron flux computation. This stems from the fact that the relic density in previous works is often approximated by pulling out $\langle\sigma v\rangle$ from the integral between $T_{\text{freeze-out}}$ and the current temperature. In general, full integration must be taken into account in the vicinity of a pole for an accurate calculation [36] or when the dark matter coannihilates [35]. Both of these cases often arise in various parts of the parameter space of dark matter models.

In the numerical analysis we fix $\rho_{\odot,\chi} = 0.18\ \text{GeV/cm}^3$ and take $\rho_{\odot,\psi} \in (0.18\text{–}0.25)\ \text{GeV/cm}^3$. We allow the neutralino relic density to lie in the range (0.035, 0.065) and find relatively good fits to the WMAP and PAMELA data for both the $U(1)_X$ and the $U(1)_X \times U(1)_C$ model with total relic density in the range $((\Omega_\chi h^2)_0 + (\Omega_\psi h^2)_0 = (0.08, 0.12)$. The analysis of Fig. 1 shows the PAMELA data and the positron flux ratio in the $U(1)_X$ model and in the $U(1)_X \times U(1)_C$ model consistent with assumed densities discussed above. In this fit the dominant contribution comes from the annihilation $\psi\bar{\psi} \to Z' \to e^+e^-$ and only small boost (clump) factors are used here, i.e., $B_{\bar{e}} = (2\text{–}5)$. Thus, the Breit-Wigner enhancement [12] plays an important role in achieving a simultaneous fit to the relic density and to the positron excess. Indeed, the annihilation near a Breit-Wigner pole gives a significant enhancement to the annihilation cross section of Dirac dark matter in the galaxy obviating the necessity of using large boost factors. At present, we are guided by the WMAP and PAMELA data on the mass splittings between the Dirac component of dark matter and the vector boson mass. Generically the required splitting is $2M_\psi - M_{Z'} \sim (10^2\text{–}10^3)$ MeV depending on the leptophilic nature of the models and narrowness of the resonance. In Fig. 1 the $U(1)_X \times U(1)_C$ has an 80 MeV mass splitting and the $U(1)_X$ model has 1300 MeV mass splitting. We note that since the $U(1)_X$

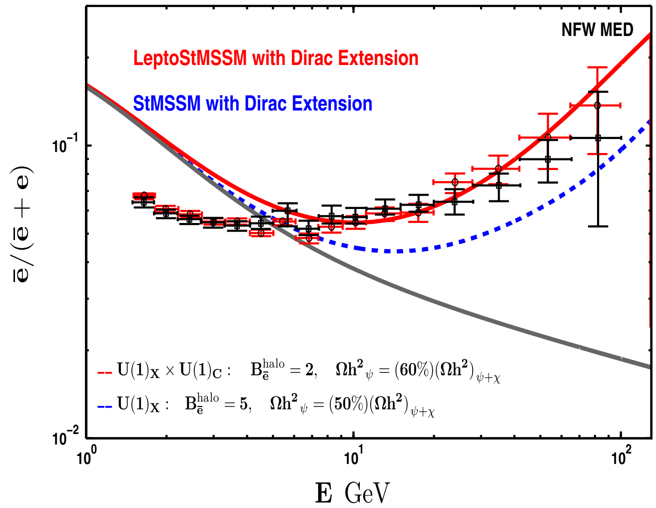

FIG. 1 (color online). Breit-Wigner enhancement [12] and the PAMELA positron excess. The analysis presented here is given in the two component Dirac-Majorana $U(1)_X$ and $U(1)_X \times U(1)_C$ models assuming a relic density decomposition of the Dirac and Majorana as given in the figure. The dominant contribution to the positron flux comes from the annihilation of the Dirac particles in the galaxy.

model is less leptophilic than the $U(1)_X \times U(1)_C$ model, we have used from the parameters of the model a smaller annihilation cross section for the $U(1)_X$ model than for the $U(1)_X \times U(1)_C$ model throughout the analysis as indicated by the positron ratio in the figure. This is motivated by the constraints from the antiproton fluxes. Thus, the $\bar{p}$ flux takes the form

$$\Phi_{\bar{p}}(T) = \frac{\eta v_{\bar{p}}}{4\pi} \frac{\rho_\odot^2}{M_{\text{DM}}^2} B_{\bar{p}} R(T) \sum_f \langle\sigma v\rangle_{f,\text{halo}} \frac{dN_{\bar{p}}^f}{dT}, \qquad (45)$$

where T is the kinetic energy, and $R(T)$ has been fit as in Ref. [40] for various profile/diffusion models and background estimates have been obtained in [46]. The antiproton flux observed at the top of the atmosphere including solar modulation can be accounted for by replacing $\Phi_{\bar{p}}(T) \to \Phi_{\bar{p}}(T + |Z|\phi_F)$ and including a kinetic energy correction ratio. We take the Fisk potential ϕ_F as 500 MV. In Fig. 2 we exhibit an analysis of the antiproton flux (signal plus background) in the $U(1)_X$ model. The analysis is done with the Navarro-Frenk-White (NFW) median model while the minimum diffusion model is unconstrained by the $\bar{p}$ data and is not shown. The PAMELA data exhibited in Fig. 2 is taken from [47]. For the $U(1)_X$ model the antiproton flux overshoots a little bit beyond $E = 10$ GeV but still lies within the limits of acceptability. For the $U(1)_X \times U(1)_C$ model, the Z' and Z'' are both leptophilic and there are no annihilations of $\psi\bar{\psi}$ into $q\bar{q}$. Because of the absence of $q\bar{q}$ final states in the annihilation, there is no contribution to the antiproton flux from the annihilation of the Dirac component of dark matter, thus

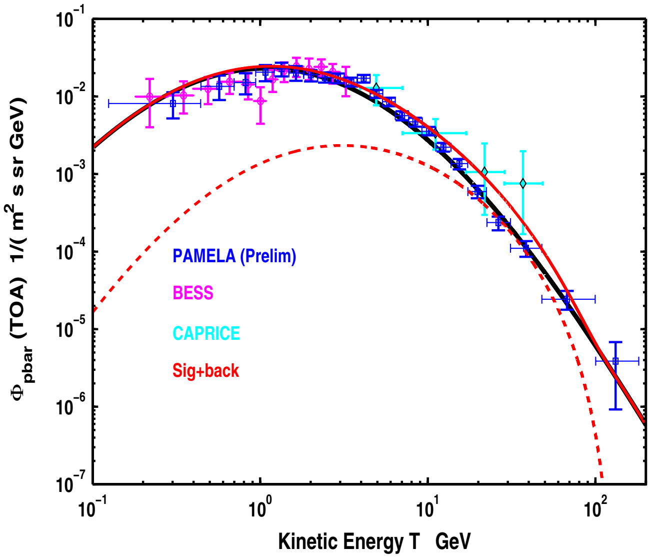

FIG. 2 (color online). Antiproton flux in the $U(1)_X$ (signal and signal plus background). The analysis is done with the NFW median model while the minimum diffusion model is unconstrained by the $\bar{p}$ data [47] and is not shown. For the $U(1)_X \times U(1)_C$ leptophilic model, the $\bar{p}$ flux from the Dirac component of dark matter does not contribute to a signal. No boost factor from clumping is taken. In both models the Majorana flux is highly suppressed relative to the Dirac flux, and is thus not shown separately. Annihilation cross sections and local dark matter densities are those used in Fig. 1.

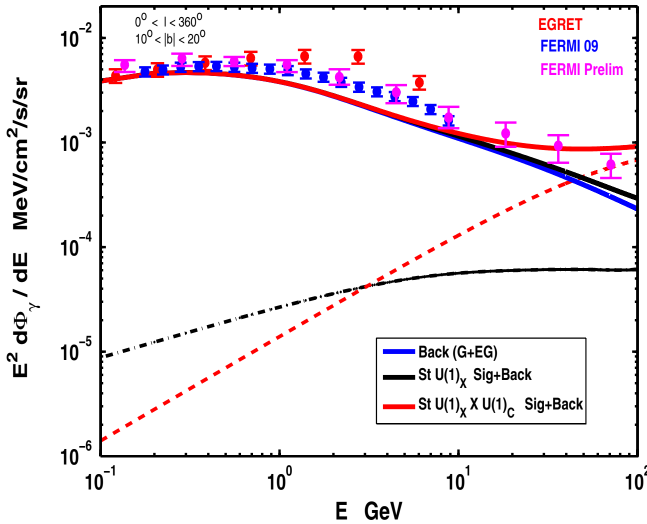

FIG. 3 (color online). Photon flux in the $U(1)_X$ and $U(1)_X \times U(1)_C$ models. The photon flux for $U(1)_X$ includes contributions from the quarks and taus and bremsstrahlung, while the photon flux for the $U(1)_X \times U(1)_C$ model is highly suppressed at low energies and peaks at larger energies from the bremsstrahlung. Annihilation cross sections and local dark matter densities are those as in Figs. 1 and 2.

the prediction of the model is not observable above the background. Finally, we look at the photon flux. In the angular region $\Delta\Omega$ the (differential) photon flux (sometime denoted $d\Phi_\gamma/dE$) is given by

$$\Phi_\gamma = \frac{\eta}{4\pi}\frac{r_\odot \rho_\odot^2}{M_{\rm DM}^2}\sum_f \langle\sigma v\rangle_{f,\rm halo}\frac{dN_\gamma^f}{dE}\bar{J}\Delta\Omega,$$
$$\bar{J} = \frac{1}{\Delta\Omega}\int_{\Delta\Omega}\int_{\rm los}\frac{ds}{r_\odot}\left(\frac{\rho(r(s,\psi))}{\rho_\odot}\right)^2. \tag{46}$$

For the $U(1)_X$ model there are three contributions to the photon flux. These arise from the $q\bar{q}$, $\tau\bar{\tau}$, and from bremsstrahlung (see i.e. [48,49]). For the $U(1)_X \times U(1)_C$ model since the Z' and Z'' are leptophilic with allowed final states being only in the first two generations of leptons, there are no final states of the type $q\bar{q}$ and $\tau\bar{\tau}$ for the Dirac component. However, there is an emission of continuum radiation for the Dirac component (see e.g., [50] using PYTHIA [51]) which arises because $\psi\bar{\psi}$ annihilate into $e\bar{e}$ and $\mu\bar{\mu}$ and there is an associated photon continuum radiation from bremsstrahlung.

In Fig. 3 we give an analysis of the continuum photon flux in the angular region where the integral over the line of sight is rather insensitive to the details of the dark matter distribution [52,53] (for a recent analysis with focus on the galactic center see [54]). The analysis is given for both $U(1)_X$ and $U(1)_X \times U(1)_C$ models using an isothermal profile. The continuum photon flux for the $U(1)_X$ model arises mostly from $q\bar{q}$ and $\tau\bar{\tau}$ at low energies while the final state radiation, i.e. $e^+e^-\gamma$ and $\mu^+\mu^-\gamma$ takes over at high energies where $E_\gamma/M_\psi \to 1$. Also shown is the EGRET [55] data and the more recent FERMI-LAT data [56] as well as the background flux in 10–20 region as estimated in the GALPROP analysis of [57]. For the $U(1)_X \times U(1)_C$ model the total photon flux is the suppressed contribution from the Majorana and the dominant Dirac source arising from the bremsstrahlung from the final states $e^+e^-\gamma$ and from $\mu^+\mu^-\gamma$. One finds that the continuum spectrum with bremsstrahlung is in accord with the current experimental data. Regarding the monochromatic photon radiation from the annihilation of dark matter, it is suppressed by $\epsilon^2 \sim 10^{-4}$ for the $U(1)_X$ model and the prediction for this model is far below the current experimental limits. For the $U(1)_X \times U(1)_C$ model there is no coupling of the hidden sector Dirac particles to the photon if the mixing with the hypercharge gauge boson vanishes. Thus, at the tree level there would be no emission of monochromatic radiation in the annihilation of dark matter in this model. In the case of a small or nonvanishing mixing with the hypercharge as in the $U(1)_X$ model, this emission is also suppressed.

V. CDMS-II AND XENON

The CDMS-II results mentioned in Sec. I raise the possibility that 1–2 dark matter events may have been seen in the CDMS-II detector, and this possibility has led to a significant theoretical activity [58]. Many analyses within supersymmetry assume the supersymmetric cold dark matter is entirely composed of neutralinos. We analyze the event rates in CDMS-II and in XENON detectors for the case when roughly only half of the dark matter is constituted of neutralinos, a situation which holds for both the $U(1)_X$ and the $U(1)_X \times U(1)_C$ models. In the analysis we use MICROMEGAS [59] and impose the electroweak symmetry breaking constraints as well as all the current experimental constraints such as on $g-2$, flavor changing neutral currents, i.e., $b \to s\gamma$ and $B_s \to \mu^+\mu^-$ branching ratio constraints, (see i.e. [60–62] for recent analyses) and require that about half the relic density as given by WMAP be given by neutralinos.

Next we note that in the direct detection experiments, the Dirac component does not give an appreciable contribution and essentially the entire contribution to the event rates arises from the Majoranas. Specifically consider the $U(1)_X$ model. Here the event rates arising from the scattering of Dirac particles from nuclear targets are suppressed by a factor of ϵ^2 relative to what one would find in the scattering of Majoranas. This is easily seen as follows: Dirac dark matter interacts with quarks in the target particles by the exchange of γ, Z, Z'. The couplings of γ, Z to Dirac dark matter are suppressed by a factor of ϵ since Dirac dark matter resides in the hidden sector. Thus, the cross sections arising from the exchange of γ, Z are suppressed by a factor of ϵ^2. Next we consider the exchange of Z'. The coupling of Z' [which is mostly a $U(1)_X$ gauge boson] with the Dirac component of dark matter is assumed to be normal size, i.e., $O(g_X) \sim O(g_2)$. However, its coupling with quarks is suppressed by a factor of ϵ. Thus the exchange of Z' also gives a scattering cross section which is suppressed by $O(\epsilon^2)$. Since $\epsilon^2 \lesssim 10^{-4}$ the Dirac component gives a negligible number of events in the direct detection experiments relative to what the neutralinos give as far as the Z' and Z'' poles are concerned. A concrete analysis of events rates for the Dirac particles in these experiments has recently been given [63] and our estimate is in accord with this analysis. The contribution from the photon pole depends on the cutoff at small angles. Further, an analysis of the event rates is subject to absorption both by the atmosphere as well as by dirt and rock in Earth before the milli charged particle gets to the detector (see, e.g., [64]).

For the $U(1)_X \times U(1)_C$ model, the Dirac particles have no interaction with the quarks, so the contribution of the Dirac particles to event rates in the direct detection experiments is absent. Thus in either case the dominant contribution to event rates in experiments for the direct detection of dark matter comes from neutralinos. In Fig. 4 we give an analysis of spin-independent cross section in NUSUGRA models for the parameter space of supergravity models with nonuniversalities in the gaugino sector so that $m_a = m_{1/2}(1 + \delta_a)$ with $(a = 1, 2, 3)$. In the analysis the

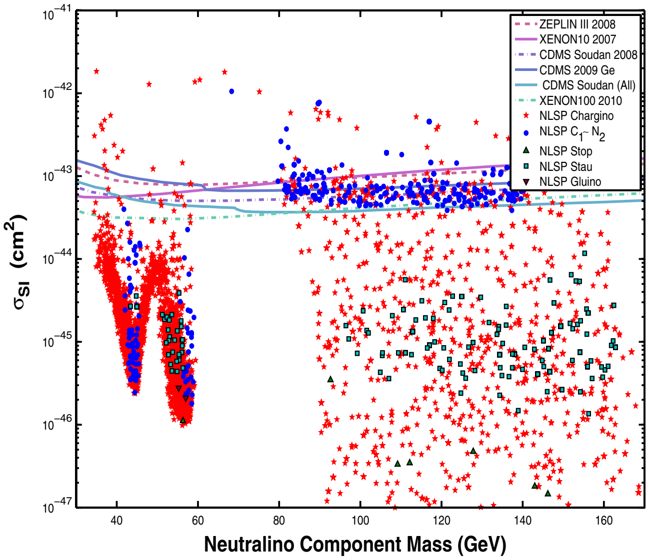

FIG. 4 (color online). An analysis of the spin-independent cross section for the parameter space of supergravity models with nonuniversalities in the gaugino sector. The analysis given above is valid for both the $U(1)_X$ and $U(1)_X \times U(1)_C$ models with a neutralino component and a Dirac component of dark matter. Models are labeled by the NLSP, which under the constraints of radiative breaking, mass limits, and flavor changing neutral currents allow chargino, stau, stop, and gluino NLSPs. By far, here the chargino NLSP arises most often.

NUSUGRA parameters are chosen in the following range: $m_0 < 3$ TeV, $m_{1/2} < 400$ GeV, $\delta_{a=2,3}$ lie in the range $(-1, 1)$, (which statistically favors the low LSP mass region which is also the region of interest in this analysis) $|A_0/m_0| < 4$, and $\tan\beta = (1\text{–}60)$. The current limits from CDMS, XENON, and from other experiments are also exhibited. We are assuming the neutralinos are contributing roughly half the relic abundance and roughly half the local density of dark matter. There is no rescaling by the dark matter density in these figures. Note the models are dominated by chargino NLSPs (next heavier beyond the LSP) [65,66] and the presence of a low mass chargino wall [67]. The relatively empty region in the range of (70–90) GeV follows from the constraint on the chargino mass being larger than 100 GeV, and an inability for the chargino and LSP to therefore coannihilate in this region, along with the constraints that the stop and stau are larger in mass than 100 GeV, and the gluino should be larger than about 300 GeV. Further, the lightest CP even Higgs has been constrained to lie higher than 110 GeV. The region of low mass is mostly controlled by the poles of the MSSM, while the higher mass region above 100 GeV is controlled mostly by coannihilations.

VI. COLLIDER SIGNALS

We discuss now the collider implications of the $U(1)_X$ and the $U(1)_X \times U(1)_C$ models. In Table I we give some

TABLE I. Top section of the table: A sample set of NUSUGRA models which produce the Majorana component of dark matter and makes up about half the relic density of the universe for both the $U(1)_X$ and the $U(1)_X \times U(1)_C$ two component models. Middle section of the table: Masses for light sparticles including the neutralino χ, the light chargino $\chi^{\pm}$, the gluino ($\tilde{g}$), the light stau ($\tilde{\tau}_1$), the light stop ($\tilde{t}_1$), and the CP even Higgses h, H/A (charged Higgs is slightly heavier) for the same set of inputs as given in the top table. Bottom section of the table: The Majorana relic density $(\Omega h^2)_\chi$, $\langle\sigma v\rangle^{\chi}_{\rm halo} \ll \langle\sigma v\rangle^{\psi}_{\rm halo}$, spin-independent neutralino proton cross section $\sigma(SI)_{\chi p}$, and event rates/kg/day in germanium and xenon detectors corresponding to entries in the top table with an assumed 30% efficiency and with $\rho_{\odot,\chi} \sim (1/2)\rho_{\odot,\rm total} \sim 0.18$ GeV/cm^3. The analysis is done with a top pole mass at 171 GeV and the models show stability in the relic density with small changes in the pole mass. The $\sigma(SI)$ given here are also within the range of XENON-100 and are on the edge of the limits recently reported in Ref. [14]. More low mass models can be seen in Fig. 4.

m_0GeV	$m_{1/2}$GeV	A_0GeV	$\delta_{1,2,3}$	$\tan\beta$
676	148	118	(− 0.26, 0.63, 0.53)	41
300	195	−100	(− 0.41, 0.43, 0.53)	41
767	173	−444	(− 0.38, 0.60, 0.09)	45
1718	297	1736	(0, −0.37, −0.68)	47
1973	227	1209	(0, −0.34, 0.34)	28
1152	139	1551	(0, 0.42, 0.03)	50
2174	314	1537	(0, 0.60, 0.18)	25
$m_{\chi_1^0=\chi}$GeV	$m_{(\tilde{\chi}^{\pm},\tilde{g})}$GeV	$m_{(\tilde{\tau}_1,\tilde{t}_1)}$GeV	m_hGeV	$m_{A\sim H}$GeV
42	(168, 595)	(556, 561)	110	480
44	(210, 723)	(211, 514)	112	370
42	(201, 515)	(578, 508)	111	438
120	(134, 311)	(1233, 943)	114	718
86	(103, 833)	(1822, 1259)	116	1666
54	(142, 418)	(708, 639)	110	475
106	(144, 982)	(2042, 1402)	116	1920
$(\Omega h^2)_\chi$ (half the DM)	$\langle\sigma v\rangle^{\chi}_{\rm halo}$cm^3/s ($\chi$ halo cross section)	$\sigma(SI)_{\chi p}$cm^2 Direct detection	Ge (evts/kg/day) ([10–50] KeV, all)	Xe (evts/kg/day) ([10–50] KeV, all)
5.1×10^{-2}	5×10^{-28}	4×10^{-44}	$(6 \times 10^{-3}, 1 \times 10^{-2})$	$(7 \times 10^{-3}, 2 \times 10^{-2})$
6.5×10^{-2}	4×10^{-28}	3×10^{-44}	$(4 \times 10^{-3}, 1 \times 10^{-2})$	$(6 \times 10^{-3}, 2 \times 10^{-2})$
5.5×10^{-2}	6×10^{-28}	4×10^{-44}	$(6 \times 10^{-3}, 1 \times 10^{-2})$	$(8 \times 10^{-3}, 3 \times 10^{-2})$
4.6×10^{-2}	2×10^{-26}	3×10^{-44}	$(4 \times 10^{-3}, 7 \times 10^{-3})$	$(5 \times 10^{-3}, 1 \times 10^{-2})$
5.1×10^{-2}	2×10^{-26}	2×10^{-44}	$(4 \times 10^{-3}, 7 \times 10^{-3})$	$(5 \times 10^{-3}, 1 \times 10^{-2})$
5.1×10^{-2}	3×10^{-28}	2×10^{-44}	$(3 \times 10^{-3}, 6 \times 10^{-3})$	$(4 \times 10^{-3}, 1 \times 10^{-2})$
5.8×10^{-2}	3×10^{-26}	5×10^{-44}	$(8 \times 10^{-3}, 1 \times 10^{-2})$	$(1 \times 10^{-2}, 2 \times 10^{-2})$

concrete models which generate half the relic abundance from the neutralino dark matter. These models produce event rates in germanium and in xenon at detectable levels with a relatively light spectrum. The predicted spin-independent elastic WIMP-nucleon cross sections are on the edge of the limits reported by XENON-100 [14]. Specifically, all of the models listed in Table I have a light neutralino mass and several also have light Higgses (for recent work relating to light Higgses and the COGENT [68] and DAMA [69] data, see [70–73]). Further, essentially all the models in the table have a gluino lying in the sub-TeV range and typically all the charginos are light. However, some of the models have rather large scalar masses in the (1.5–2) TeV region indicating that they originate on the hyperbolic branch of radiative breaking of the electroweak symmetry [74]. The models listed in Table I share the property that the gauginos in all cases are relatively light. Thus, such models should give rise to detectable signals in the form of leptons and jets and missing energy at the LHC with modest luminosity (though the missing energy may be difficult to estimate in early runs). The models with very light gluinos could surface with less than 1 fb^{-1} at LHC center of mass energies $\sqrt{s} =$ 7, 10 TeV, (for recent analyses see [19,75,76]) while many of the models should be discoverable with $O(10)$ fb^{-1} at $\sqrt{s} = 10$ TeV. (In fact one can glean this from the analysis of the first listing of Ref. [58].) As many of the candidate models have rather light gluinos with a chargino NLSP, such models likely will produce missing energy which is very SM like. Large event rates can arise, however, from multijets, and, in particular, from b jets. One also expects a sizable amount of leptons in these models. For the cases where stau-coannihilation survives, the leptonic signals are likely to be stronger and the missing energy larger than for the chargino NLSP cases. However, since these models have very low SUSY scales, most of them should indeed be discoverable (and likely rather early) at the LHC. Additionally for the $U(1)_X \times U(1)_C$ model Z'' offers the possibility of discovery at a next linear collider (NLC) as it will have distinct signatures. Its decay width is significantly smaller than what a GUT-type Z' with the same mass will have. Additionally, it has visible decays only into $e\bar{e}$ and $\mu\bar{\mu}$ along with radiation at the NLC. Thus, a Z'' of this type can be detectable at an NLC because of its distinct signatures. However, a full simulation of collider signals requires a separate dedicated analysis.

VII. CONCLUSION

In this work we have proposed a new class of models with dark matter consisting of two, three, or even four components. We considered the two component model consisting of Dirac and Majorana particles in detail. We showed that this two component model can fit the positron excess seen in the PAMELA experiment as well as can produce detectable signals in the current direct detection experiments while satisfying WMAP relic density constraints. Thus, the Dirac component of the two component dark matter model allows a fit to the PAMELA data via annihilation of the Dirac particles close to a Breit-Wigner pole. On the other hand, the Majorana component of dark matter plays the dominant role in the generation of events in dark matter detectors. Specifically, we showed that in the two component picture it is possible to generate events of size 1–2 in 612 Kg-d of data in the CDMS-II detector as well as event rates that can be tested by the results of XENON-100, and an observable number of events in other ongoing direct detection experiments. Further, it was shown that models which lead to detectable signals in direct detection experiments are typically associated with a relatively light spectrum which is discoverable at the LHC with modest luminosity. Further, one class of models discussed in this work produces a Z'' vector boson which has visible decays only to e^+e^- and $\mu^+\mu^-$. The proposed models contains massive scalar fields which are also possible candidates for dark matter. Thus these spin zero fields in combinations with Dirac and Majorana particles present the possibility of a multicomponent dark matter. Finally, we note that it would be very interesting to investigate phenomena where both components play a significant role, i.e., regions of the parameter space of the model which allow both dark matter candidates to appear in the data analysis. Such a possibility may appear in certain decay fragments at the LHC where the missing energy signals from the two dark matter particles would be different because of their different masses and interactions.

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