Breaking the dark force

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ABSTRACT: Recently Arkani-Hamed, Finkbeiner, Slatyer and Weiner proposed a unified explanation to the set of experimental and observational anomalies with possible connections to Dark Matter (DM). A central role is played by GeV scale "dark" gauge bosons exchanged by the weak scale DM. Motivated by this proposal, we build an explicit model of DM in the context of Weak Scale Supersymmetry (SUSY). We employ high-scale SUSY-breaking and invoke the Giudice-Masiero mechanism to generate the weak scale DM masses. By sequestering the dark sector from the SUSY-breaking "hidden" sector, it naturally acquires GeV scale soft masses that help generate the dark gauge boson mass. The visible MSSM sector is not fully sequestered and acquires gaugino-mediated soft terms at the weak scale. This hierarchy of scales naturally leads to the Sommerfeld enhancement of DM annihilations needed to account for the e^+e^- excesses in the PAMELA and ATIC experiments. The possibility of co-existing species of DM is used to show how the INTEGRAL and DAMA anomalies can both be explained. We study the cosmological constraints on the new stable or long-lived light particles that appear in our models. We discuss the lepton-jet collider signals suggested by Arkani-Hamed and Weiner, and find that they are not guaranteed in our construction.

Contents

1 Introduction

In recent years there have been an increasing number of experimental and observational anomalies which might possibly be due to the effects of Dark Matter (DM). INTEGRAL [1] has detected a 511 keV emission line from the center of the galaxy due to positron annihilations, DAMA [2] found annual modulations in their direct detection experiment, PAMELA [3] measured a positron excess in the high-energy e^+e^- flux, while ATIC [4] measured an excess in the overall e^+e^- flux at even higher energies. While any or all of these experiments might ultimately be explained without recourse to new physics, it is obviously of interest to tentatively consider them as clues to the nature of DM, especially if such considerations motivate the search for new types of experimental signals.

If the DM explanation for DAMA and INTEGRAL experiments are accepted, one should consider non-minimal DM scenarios. It is not easy to reconcile the DAMA measurements with other direct detection experiments. One of the notable explanations to this disagreement is the idea of inelastic dark matter (iDM) [5]. On the other hand, the 511 keV line seen by INTEGRAL can be successfully explained by the scenario of exciting

dark matter (XDM) [6, 7]. It is striking that these two phenomenologically distinct anomalies can be explained by similar mechanisms, namely roughly $\mathcal{O}(\text{MeV})$ splittings between DM states.[1]

The explanation of PAMELA/ATIC signals via DM annihilation is also challenging: one requires DM annihilation cross section two to three orders of magnitude larger than needed to account for their measured relic abundance [10]. Introducing such large boost factors is somewhat implausible. Moreover, the absence of observable antiproton [11] or gamma-ray excesses[2] imposes stringent constraints on any explanation in particular requiring suppression of hadron production from DM annihilations.

Recently a single attractive picture, explaining simultaneously all these phenomena, was suggested in [16, 17]. The main idea is that the DM is charged under some intermediate-range gauge force with a gauge boson mass at the GeV scale, discussed earlier in [6, 18]. This would result in a significant Sommerfeld enhancement of the late-time DM annihilation cross section, once DM is non-relativistic. This framework also explains why there is an excess of positrons but no excess in the antiproton or photon fluxes. The main annihilation channel of DM is into the dark sector $\sim$ GeV-mass gauge bosons. The dominant interaction of the dark sector with the SM is naturally due to kinetic mixing between the dark gauge boson and the regular photon (more fundamentally, hypercharge) [19]. If these gauge bosons do not have efficient decay channels within the dark sector, they will decay via this mixing to sub-GeV charged SM particles, leading mostly to decays to e^+e^- and with few hard photons and no antiprotons. The presence of such a light subsector interacting with DM can also explain the small splittings between the DM states, required for iDM and XDM. These splittings can be either radiatively induced, or come from couplings between DM and the light fields.

Inspired by the suggestion of [16] we develop in this paper a concrete model of the dark sector, within the paradigm of Weak Scale Supersymmetry. Quite apart from the experimental anomalies discussed above, we are also motivated by trying to understand what simple connections might exist within SUSY field theory between our visible sector and other self-interacting "dark" sectors without SM-charges, in keeping with Hidden Valley ideas [20].

The central issue for the DM modelling is to find natural origins for the two scales of the dark sector, the DM mass scale of hundreds of GeV, and the GeV mass scale of the dark gauge boson. Clearly, we would like to connect both scales, as well as the weak scale, to the scale of SUSY breaking. We follow closely the idea suggested in [16] that DM gets its mass from the same mechanism responsible for the appearance of the $\mu/B\mu$ terms in the MSSM, while the GeV scale emerges radiatively with one-loop suppression, and MeV DM splitting emerge radiatively with two-loop suppression.

What is new here is that we take the simplest, and in many ways the most attractive, approach to generating a μ term and DM mass, namely (generalizations [21] of) the Giudice-

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[1]For further discussion of some of the subtleties of these ideas see [8, 9].

[2]The most stringent constraints on the DM models explaining PAMELA are imposed by HESS observations of the Galactic Center [12] and the Galactic Ridge [13]. See [14, 15] for an exhaustive discussion of the relevant constraints.

Masiero mechanism [22].[3] This mechanism works best in the context of high-scale SUSY-breaking. We will therefore consider the MSSM soft terms to be generated by gaugino-mediated SUSY breaking [24, 25], so as to resolve the SUSY flavor problem. But we are then faced with the question of why the dark gauge boson has mass of only GeV. If one introduces dark Higgs fields to generate a mass, they should naively feel gravity-mediated soft masses of order $m_{\frac{3}{2}} \sim 100$ GeV, making a GeV scale Higgs mechanism quite unnatural. We resolve this puzzle in a simple way, by taking the entire dark sector to be sequestered from the SUSY-breaking hidden sector. Our set-up can therefore be pictured as in figure 1.

The sequestered dark sector fields get their masses from several different sources, but all of them are naturally suppressed compared to the scale $m_{\frac{3}{2}}$. First, the supersymmetric gauge kinetic mixing also contains D-term mixing, resulting in an effective Fayet-Iliopoulos (FI) term for the dark sector after electroweak symmetry breaking [26]. The strength of this mixing, ϵ, can naturally be of one-loop to two-loop order if it is due to loops of heavy "link fields", carrying both dark charge and hypercharge (and possibly other SM gauge quantum numbers). With this suppression, the effective FI term in the dark sector is of order GeV^2. Furthermore, the sequestered dark sector also gets anomaly-mediated (AMSB) soft masses [21, 27]. These masses are generically of order $m_{\frac{3}{2}}/(16\pi^2) \sim 1$ GeV. Moreover, IR-free gauge forces give tachyonic contributions to charged scalar masses in AMSB. These contributions combined with the effective FI term will break the dark force at $\sim$ GeV, as required.

Careful analysis of the dark sector in our construction reveals several interesting features, which seem quite generic. We emphasize that almost all the models of this type involve *co-existing DM*. Besides the stable heavy DM (HDM), which is responsible for the experimental signals listed above, one finds a stable lightest R-odd particle at $\sim$ GeV (of course, if R-parity is assumed). This particle belongs to the dark sector and forms a light DM (LDM). In our model, the LDM is the lightest fermion of the dark sector, which is thermally produced. This imposes a stringent constraint on the parameter space, to ensure that the LDM has sufficiently strong annihilation channels, such that its relic abundance is small.

The dark sector can also contain long-lived particles. Indeed, if the dark gauge boson is not the lightest particle in the dark spectrum, the lightest particle decays induced purely by gauge kinetic mixing can be extremely slow. Such long-lived particles are severely constrained by Big Bang nucleosynthesis (BBN). However we identify simple new couplings between the dark and visible sectors that greatly speed up the decays. We also discuss how these couplings and the particle spectrum affect the "lepton-jets" signatures predicted in [17]. We stress that these collider signals are not an inevitable outcome of our model. Dark particles produced at colliders can be sufficiently long-lived to escape the detectors before their decays to leptons.

Finally, in order to explain both DAMA and INTEGRAL by iDM and XDM respectively, we are required to introduce two flavors of HDM. Because the HDM is charged

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[3]For an earlier attempt to realize the model suggested in [17] in the framework of gauge-mediation, see [23].

Figure 1. A qualitative picture of extra-dimensional sequestering.

only under an Abelian symmetry in our model, we do not have enough degeneracy in a single DM flavor to have both iDM and XDM splittings. While such co-existing HDM is a plausible and interesting possibility, we also note that our model might also be extendable to accommodate non-Abelian dark gauge group, enabling a simultaneous explanation for iDM an XDM from a single multiplet of DM.

The paper is organized as follows. In the section 2, we assemble our model in a modular fashion. In section 3, we derive the spectrum of the dark sector taking into account SUSY breaking. In section 4, we discuss phenomenological requirements of our model and demonstrate viable regions in the parameter space. In particular we explain how we get the required type of experimental signals for PAMELA/ATIC. We discuss the possible presence or absence of lepton jets in collider signals, and its sensitivity to the flavor structure of our new couplings between the dark and visible sectors. Section 5 provides our conclusions.

While this paper was in preparation we learned about the work of [28], which has a considerable overlap with our work. While both our models consider a similar field content for the dark sector, the implementation of SUSY breaking is different.

2 The model

We consider high-scale supersymmetry breaking, gaugino-mediated to the visible sector, so that the mass scale of the SM superpartners is of order of the gravitino mass $m_{\frac{3}{2}}$. This requires sequestering of MSSM matter, but not gauge fields, from the SUSY breaking hidden sector. We also assume that the *entire* dark sector and any link fields are sequestered from the SUSY breaking hidden sector. The basic set-up is then shown in figure 1.

Our dark sector is a simple generalization of SQED, with gauge group $U(1)_D$. We choose some massive vector-like charged pairs, $X_i + \bar{X}_i$, to house dark matter, and a massless vector-like charged pair, $T + \bar{T}$, to house the dark Higgs field responsible for breaking the $U(1)_D$ and giving mass to the dark photon. For simplicity, all the charges are taken as ± 1. We also assign $\mathbb{Z}_2$ symmetries under which the $X_i, \bar{X}_i$ are odd, to ensure

their stability (and non-mixing) when other interactions are added. All the chiral fields are taken R-parity even. None of them carry SM quantum numbers.

After $U(1)_D$ is broken, the safest option is to have the entire dark sector acquire mass. If there are massless particles in the dark sector, then dark matter annihilations will predominantly end up in these particles, and there is no chance of decaying into the SM to give the PAMELA/ATIC signals. From this point of view, we can anticipate a general problem if we look at the fermions in $T, \bar{T}$. Without knowing the details of soft SUSY breaking in the dark sector we can still conclude that one linear combination of these fermions will mix with the dark gaugino after the Higgs mechanism, giving mass to this combination in general. But the remaining linear combination of charged fermions remains massless, protected by the chiral symmetry of the massless fields. Of course we could add a small mass term for $T, \bar{T}$, but this would then require a tuned value, posing a kind of dark "μ-problem". Instead we will introduce a gauge singlet, S, whose fermion will marry the otherwise-massless charged fermion to form a Dirac state, once we include a superpotential term

$$W = \lambda S T \bar{T} + \frac{\kappa}{3!} S^3. \tag{2.1}$$

Of course S is also taken as even under all discrete symmetries.

We consider the dark gauge coupling, g_D, as well as λ, κ to be roughly order one.

2.1 Dark matter mass

We do not put in direct superpotential mass terms in the dark sector, because just as for the MSSM Higgs multiplet, this would lead to a μ problem, or coincidence of scales arising from disparate sources. Of course, since such mass terms are not renormalized and are protected by chiral symmetries, their vanishing below the Planck scale is a natural possibility.

We are assuming that the entire dark sector is sequestered from the hidden sector responsible for intermediate scale supersymmetry breaking, so we use a simple generalization of the Giudice-Masiero mechanism [21] to give weak-scale masses to the X's. Their Kahler potential is given in off-shell SUGRA by

$$K = |\phi|^2 \left(|X|^2 + |\bar{X}|^2 + c_i (X_i \bar{X}_i + \text{c.c.}) \right), \tag{2.2}$$

where the VEV of the SUGRA conformal compensator is given by

$$\phi \equiv 1 + m_{\frac{3}{2}} \theta^2. \tag{2.3}$$

We take the natural size of dimensionless coefficients $c_i \sim \mathcal{O}(1)$. One can think of them as arising at the Planck scale in the form $\mathcal{O}(1) \langle \chi^\dagger \rangle / M_{\text{Pl}}$, where χ is a chiral field with $\sim$ Planckian mass and VEV that breaks the chiral symmetries that protect a direct X mass.

After superfield rescaling $X\phi \to \phi$ and $\bar{X}\phi \to \phi$ we get the effective Kähler potential

$$K = |X|^2 + |\bar{X}|^2 + c \frac{\phi^\dagger}{\phi} (X\bar{X} + \text{c.c.}) \tag{2.4}$$

resulting in μ and $B\mu$ terms

$$\mu_i = c_i m_{\frac{3}{2}}, \quad B_i \mu_i = c_i m_{\frac{3}{2}}^2. \tag{2.5}$$

Assuming that the MSSM acquires SUSY breaking through gravity-mediation or gaugino-mediation implies that $m_{\frac{3}{2}}$, and therefore μ_i, B_i, are of order the weak scale. Note that these "soft" masses have emerged despite sequestering of the dark fields (no direct coupling to the hidden sector).

The tree-level spectrum of $X + \bar{X}$ is given by

$$
\begin{aligned}
m_{1/2} &= c m_{\frac{3}{2}} \\
m_{0,X+\bar{X}^\dagger}^2 &= c(c+1) m_{\frac{3}{2}}{}^2 \\
m_{0,X-\bar{X}^\dagger}^2 &= c(c-1) m_{\frac{3}{2}}{}^2.
\end{aligned}
\tag{2.6}
$$

We take $|c_i| > 1$ to avoid scalar tachyons. We take the analogous c coefficient to vanish for our dark Higgs sector, $T + \bar{T}$, to avoid getting mass. Again, this is a natural possibility.

2.2 Mixing with SM hypercharge

The renormalizable kinetic mixing of $U(1)_D$ and hypercharge is given by

$$
\mathcal{L} \supset \frac{\epsilon}{2} \int d^2\theta W_D^\alpha W_{\alpha Y} + \text{c.c.}.
\tag{2.7}
$$

We assume that ϵ is induced by loops of heavy "link fields" carrying both SM and dark charges, with a natural range of $\epsilon \sim 10^{-4} - 10^{-3}$.

For the most part such a tiny mixing has little effect on the dark sector, but there are three qualitatively important exceptions which we can anticipate:

(i) As suggested by [16] otherwise stable particles in the dark sector may decay into the SM via this tiny mixing. This is the key to the PAMELA/ATIC signals from dark matter annihilations in their proposal.

(ii) The mixing allows photon-dark-photon exchanges to mediate direct detection of dark matter. The smallness of the mixing will be offset by the lightness of the mediating particle relative to say Z-exchange in more standard scenarios.

(iii) As noted in [26], the D-terms of the U(1)'s are also mixed, and electroweak symmetry breaking in the SM leads to an effective Fayet-Iliopoulos D-term in the dark sector:

$$
\xi_{\text{FI}} = \epsilon D_Y, \quad D_Y = \frac{g' v^2 \cos(2\beta)}{2}
\tag{2.8}
$$

where D_Y denotes the D-term of the SM hypercharge and v is the electroweak scale 176 GeV. For the range of ϵ considered this introduces the desired GeV scale into the dark sector.

Note that the decays to the SM via (i) are most efficient if the dark photon is stable for $\epsilon = 0$, since it is directly coupled to the SM photon by the mixing. In the present paper we will consider this to be a subdominant channel for decays to the SM, to satisfy phenomenological constraints discussed in section 4. In the next subsection we will introduce other simple couplings between the two sectors which will dominate the PAMELA/ATIC signals from dark matter annihilations. The kinetic mixing ϵ will however continue to serve purposes (ii) and (iii).

2.3 Singlet couplings to the SM

We will be led to study a part of the parameter space in which the singlet scalar is the lightest particle, into which other dark gauge/Higgs degrees of freedom rapidly decay or annihilate. The PAMELA/ATIC signals can then arise from decays of the singlet scalar into the SM. Rapid decays can arise from non-renormalizable direct couplings of the form

$$\mathcal{L} \supset \int d^2\theta \left(\frac{SLH_d\bar{e}}{\Lambda} + \frac{S\bar{d}H_dQ}{\Lambda} \right), \tag{2.9}$$

where $L, \bar{e}, Q, \bar{d}$ denote the usual weak doublets and down-type singlets of the MSSM and H_d is the down-type Higgs. We might expect the coefficients of these couplings to be roughly similar to Yukawa couplings, while the overall suppression scale Λ is due to integrating out some high-scale physics.

A simple example of such high-scale physics is a vector-like pair of chiral multiplets carrying SM quantum numbers with Yukawa couplings to the SM and S. To preserve perturbative unification we can take the pair to consist of a massive lepton doublet[4], $\mathscr{L} + \bar{\mathscr{L}}$, and massive down-type quark singlet, $\mathscr{D} + \bar{\mathscr{D}}$, forming a $5 + \bar{5}$ of SU(5). Given superpotential couplings of the form

$$W = SL\bar{\mathscr{L}} + \bar{\mathscr{L}}H_d\bar{e} + S\bar{d}\mathscr{D} + \bar{\mathscr{D}}H_dQ + M_{\mathscr{L}}\mathscr{L}\bar{\mathscr{L}} + M_{\mathscr{D}}\mathscr{D}\bar{\mathscr{D}}, \tag{2.10}$$

we get (2.9) upon integrating out the mass $M_{\mathscr{L}}, M_{\mathscr{D}}$ states.

Upon electroweak symmetry breaking, the effective couplings (2.9) give the scalar S decay channels into kinematically available pairs of SM leptons and quarks. There is a large range of possible masses M or scales Λ over which these decays can be very rapid. We will return to discuss the constraints on these decays and their phenomenology in section 4.

2.4 Dark matter splittings for iDM and/or XDM

Our stable dark matter particles are the complex dark-charged scalars $X_i - \bar{X}_i^\dagger$. For iDM or XDM to work we need to split the real and imaginary components of these scalars which are connected by dark-photon couplings, by of order MeV. If we want *both* iDM and XDM we will need sub-MeV and super-MeV splittings, and therefore we will simply invoke two co-existent dark matter species, $i = 1, 2$. If only one or the other of iDM or XDM is needed, one species of dark matter, $i = 1$, will suffice.

Real-imaginary splittings can emerge once the dark matter feels the breaking of $U(1)_D$, say via the coupling to the dark Higgs,

$$\mathcal{L} \supset \mathcal{O}(1)\frac{\int d^2\theta T^2\bar{X}^2}{\text{TeV}} \tag{2.11}$$

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[4]Somewhat similar idea of introducing new leptons was suggested in [29]. Note that in our case these fields can be much heaver than ~ 100 GeV as was proposed in [29].

Such non-renormalizable couplings can readily arise by integrating out $\mathbb{Z}_2$-odd singlets, A_i, with weak scale masses via the (conformal compensator version of the) Giudice-Masiero mechanism, and couplings of the form

$$\mathcal{L} \supset \eta \int d^2\theta \, T\bar{X}_i A_i \,, \tag{2.12}$$

where we define a coefficient η for later reference.

Now we have introduced a full theory of the dark sector. Adding the (gauged) kinetic terms for all the fields, together with the equations (2.1), (2.4), (2.7), (2.9), (2.11) one gets the full effective Lagrangian of the dark sector and its interactions with the visible sector, to be used in the analysis below.

3 Dark spectrum

In this section we will determine the ground state of the dark sector in various parts of the parameter space as a consequence of the SUSY breaking. We will then study higgsing of $\mathrm{U}(1)_D$ and the mass spectrum.

3.1 SUSY breaking contributions to the dark sector

SUSY breaking enters the dark sector from three different sources:

- Effective Fayet-Iliopoulos term

- standard AMSB

- deflection from AMSB due to non-decoupling threshold effects

Below we discuss each of these contributions.

As discussed in the previous section gauge kinetic mixings, electro-weak and SUSY breaking in the visible sector induce an effective FI term

$$\mathcal{L} \supset \xi_{\mathrm{FI}} \int d^4\theta V \tag{3.1}$$

for the dark sector (2.8). As we have already mentioned for the values of ϵ of order 10^{-3} – 10^{-4}, $\xi_{\mathrm{FI}} \sim (1 \text{ GeV})^2$. This contribution is the only non-sequestered one in the dark sector. (Sequestering would be exact in the $\epsilon \to 0$ limit.) While it is triggered by SUSY breaking effects in the visible sector, from the dark sector point of view this contribution is entirely supersymmetric (3.1).

Because of sequestering the remaining SUSY breaking contributions must arise from AMSB [21, 27]. We are interested primarily in soft terms for the light degrees of freedom. The usual form of such contributions depend only on the IR effective theory and $m_{\frac{3}{2}}$. These contributions are of order one-loop for the gaugino masses and A-terms and of order two-loops for the scalar masses squared, corresponding to the $\mathcal{O}(1 \text{ GeV})$ scale again, given $m_{\frac{3}{2}} \sim \mathcal{O}(100 \text{ GeV})$.

In AMSB, heavy thresholds due to a supersymmetric mass decouple from IR soft terms. Naively, such decoupling would apply to the HDM contributions to IR soft terms. However this is not true for two reasons. Firstly, such decoupling is only true when the thresholds are much larger than $m_{\frac{3}{2}}$, while our HDM is comparable. Secondly, mass thresholds due to mass terms in a superpotential satisfy

$$\frac{B\mu}{\mu} = m_{\frac{3}{2}}, \tag{3.2}$$

while mass thresholds from the Kahler potential such as (2.4) have a sign difference

$$\frac{B\mu}{\mu} = -m_{\frac{3}{2}}. \tag{3.3}$$

While this sign difference is important for gauginos, it cancels in the dominant scalar soft terms [30].

3.2 Mass spectrum

Now we are ready to write down the scalar potential of the light dark sector. A *supersymmetric* D- and F-term scalar potential, including the effective FI term, reads

$$V = \frac{g_D^2}{8} \left(|T|^2 - |\bar{T}|^2 - \frac{\xi_{\mathrm{FI}}}{g} \right)^2 + \left| \frac{\kappa}{2} S^2 + \lambda T\bar{T} \right| + \lambda^2 \left(|TS|^2 + |\bar{T}S|^2 \right) . \tag{3.4}$$

Soft masses are

$$m_T^2 = \frac{m_{\frac{3}{2}}^2}{(16\pi^2)^2} \left(-4g_D^4 + 3|\lambda|^4 - 4|\lambda|^2 g_D^2 + \frac{|\lambda|^2|\kappa|^2}{2} \right) - \frac{1}{6} \frac{\eta^2}{16\pi^2} \left(\frac{m_{\frac{3}{2}}^2}{c^2} + \frac{m_{\frac{3}{2}}}{c_A^2} \right) \tag{3.5}$$

$$m_S^2 = \frac{m_{\frac{3}{2}}^2}{(16\pi^2)^2} \left(3|\lambda|^4 - 4|\lambda|^2 g_D^2 + 2|\lambda|^2|\kappa|^2 + \frac{3}{4}|\kappa|^4 \right) . \tag{3.6}$$

An A-term which couples the scalars $T, \bar{T}, S$ is also an important contribution to the spectrum because T condenses:

$$a \equiv a^{T\bar{T}S} = -\frac{m_{\frac{3}{2}}}{16\pi^2} \left(3\lambda|\lambda|^2 - 4\lambda g_D^2 + \frac{\lambda|\kappa|^2}{2} \right) \tag{3.7}$$

The first term in the expression (3.5) is a standard AMSB contribution while the second one comes from Yukawa couplings to HDM (2.12) with non-decoupling masses [31]. We also get soft gaugino mass, which reads (including the non-decoupling threshold (2.4))

$$m_{1/2} = \frac{g_D^2 m_{\frac{3}{2}}(N-1)}{8\pi^2} . \tag{3.8}$$

Here N denoted a number of charged HDM flavors.

First, in order to simplify the following analysis we consider $c \sim 10$ so that the non-decoupling Yukawa mediated contribution (η^2 contribution to soft masses) is subdominant. Clearly for $\lambda \gg g_D$ neither T nor $\bar{T}$ condense, so we take $\lambda < g_D$. The masses from AMSB

Figure 2. The mass spectrum of the dark sector, measured in units of GeV, plotted against various values of λ. The values which have been chosen for this plots are: $g_D = 0.4$, $m_{\frac{3}{2}} = 110$ GeV, $\kappa = 1$, $\epsilon = 10^{-4}$, $\tan\beta = 10$. Here, as well as on the following plots, the red (dotted) and the green (dashed) lines denote the scalars from S and $\bar{T}$ admixtures, the blue (dashed-dotted) line denotes a Dirac fermion from S and $\bar{T}$. A cyan solid thick line denotes the mass of the gauge boson, while the brown and the black thin lines denote Majorana fermions from the gauge multiplet and T.

are the same for T and $\bar{T}$, but the contributions of the FI term have opposite signs. For simplicity we will restrict ourselves to the regime, where only T condenses. We stress that there is no strong physical reason for restricting to this regime. Indeed even for $\epsilon = 0$ and hence vanishing FI term, we can achieve viable $U(1)_D$ breaking when all three fields T, $\bar{T}$, S get VEVs.

If only one field gets a VEV, its precise value is

$$\langle T \rangle = \sqrt{\frac{\xi_{\mathrm{FI}}}{g_D} - \frac{4m_T^2}{g_D^2}} \ . \tag{3.9}$$

Now we turn to the spectrum of the dark sector. Let us first analyze the masses of the fields in the gauge supermultiplet. When $U(1)_D$ is broken, the gauge field eats the phase of T. The mass of the gauge field is

$$m = g_D \langle T \rangle \ . \tag{3.10}$$

If the sector had been supersymmetric, this also would be a mass of the fermions. The photino mixes only with the T-fermion after T-condensation and one can write down 2×2 mass matrix for these fermions:

$$m = \begin{pmatrix} m_{1/2} & g_D \langle T \rangle \\ g_D \langle T \rangle & 0 \end{pmatrix} \tag{3.11}$$

Figure 3. The mass spectrum of the dark sector. The values chosen for the right-hand side picture are $g_D = 0.5$, $m_{\frac{3}{2}} = 100$ GeV, $\kappa = 0.8$, $\epsilon = 10^{-3}$, $\tan\beta = 8$ and for the left-hand side picture $g_D = 0.6$, $m_{\frac{3}{2}} = 85$ GeV, $\kappa = 2$, $\epsilon = 5 \times 10^{-4}$, $\tan\beta = 30$. Again the mass is measured in units of GeV, the red dotted line and the green dashed line denote scalars from S and $\bar{T}$ admixtures, the blue dased-dotted line denotes a Dirac fermion from S and $\bar{T}$. A cyan solid thick line denotes the mass of the gauge boson, while the brown and the black thin lines denote Majorana fermions from the gauge multiplet and T.

Diagonalizing this matrix we find two Majorana states with the masses

$$m = \frac{1}{2} \left| m_{1/2} \pm \sqrt{g_D^2 |T|^2 + m_{1/2}} \right| . \tag{3.12}$$

The fermion from the S multiplet joins the fermion from the $\bar{T}$ multiplet in order to form a Dirac state with the mass λT.

The scalars of S and $\bar{T}$ mix with each other by virtue of the A-term, but they do not mix with the physical scalar T. Therefore their masses squared are given by the eigenvalues of the following 2×2 complex matrix

$$m^2 = \begin{pmatrix} |\lambda|^2 |T|^2 + m_T^2 & aT \\ a^* T^* & |\lambda|^2 |T|^2 + m_S^2 \end{pmatrix} \tag{3.13}$$

The eigenvalues of this matrix are

$$m^2 = \frac{1}{2} \left(2|\lambda|^2 |T|^2 + m_S^2 + m_T^2 \pm \sqrt{(m_T^2 - m_S^2)^2 + 4|a|^2 |T|^2} \right) \tag{3.14}$$

Mostly in the parameter space, both these mass-squared are positive. Rendering one of these masses tachyonic signals the breakdown of the single-VEV solution.

We plot the spectra for various choices of parameters in figures 2 and 3. Note that for each choice of the parameters, for sufficiently small λ the lightest scalar becomes massless and below this value becomes tachyonic, signalling the onset of multi-VEV solutions, which we do not consider. In the vast part of the parameter space a scalar admixture of S and $\bar{T}$ is the lightest particle, although there is a region in which the lightest particle is one of the Majorana fermions. As we will see in the next section, the first possibility is much more preferable.

4 Phenomenology

Below we list the phenomenological issues that constrain our parameter space.

- Dark matter abundance. We should ensure that the correct dark matter density of the Universe can be accounted for as the thermal relic abundance of HDM.

- PAMELA/ATIC signal. We should provide an effective leptonic decay channel for one of the light particles in the dark sector. If the gauge boson is not the lightest particle in the spectrum, this feature is not, a-priori, guaranteed.

- LDM relic abundance. As we have pointed out, at least one state in the dark sector is absolutely stable since it is the lightest R-odd particle. One should worry that this relic is at least not overproduced. Ideally we would like to render the abundance of this relic negligible.

- Constraints from BBN. In some parts of the parameter space the lightest *unstable* particle is not the dark photon. If this particle is too long-lived, one finds severe constraints from the observed light elements abundance [32]. We should verify that our model is not excluded due to these constraints.

4.1 Who should be the lightest particle?

Observing the spectra of the dark sector (see, for example, figure 3) we distinguish between two qualitatively different cases. For relatively small λ we find that a scalar admixture of S and $\bar{T}$ is the lightest particle. The next to lightest particle is a Dirac fermion. When λ increases one of the Majorana fermions becomes the lightest state and the gauge boson is the next-to lightest state.

The lightest fermion of the hidden sector is always stable, because it is the lightest R-odd particle. We must verify that the abundance of this particle in the Early Universe is negligible. When scalars are the lightest dark particles there is no problem with the relic fermion abundance. The annihilation cross section of fermions to scalars is of order:

$$\langle \sigma v \rangle \sim \frac{\alpha_\lambda^2 |\lambda \langle T \rangle|^2}{m_T^4} \tag{4.1}$$

It is clear that this cross section is large enough for the fermions to annihilate away to scalars, and we will ensure that the scalars decay rapidly. Therefore, we do not have any problems with the LDM overproduction.

But if the Majorana fermion is the lightest particle we have a problem to wash it out efficiently enough. The only kinematically allowed annihilation channels are into the SM particles, which are strongly suppressed by ϵ, and the annihilations into the leptons are also p-wave suppressed.

We conclude that the small λ region is phenomenologically safest, where scalars are the lightest dark particles.

4.2 Decays into the SM: astrophysical implications

We will call light dark particles that cannot decay within the dark sector, "dark-stable". Such particles are either absolutely stable, forming LDM, or can decay to the SM via the very weak dark-visible couplings (2.7), (2.9). They are important for two reasons. Given that HDM can annihilate into any of the dark-stable particles, the dark-stable decays to the SM can account for a variety of astrophysical signals. At colliders, R-parity ensures that sparticle production ends in decays into the LDM, potentially with associated production of other dark particles. If these latter particles are dark-stable they can decay back to the SM, possibly within the detector.

The dark photon is dark-stable in some parts of parameter space where it cannot kinematically decay into lighter particles, as can be seen in figure 2. In this case, it nevertheless promptly decays to the SM via gauge kinetic mixing, as discussed in [17]. However, note that the dark photon can easily be heavier than a GeV, which allows it to have both leptonic and hadronic decay modes (eg. $\omega^0 + \pi^0$). As emphasized in [14], if this channel accounts for the entire PAMELA/ATIC lepton signals, then it should also produce enough photons from hadron decays to be in conflict with HESS measurements. But in our model, this is not the primary source of the PAMELA/ATIC leptons.

Indeed, the lightest dark particle in our chosen parameter region is a dark scalar, which is automatically dark-stable. Most of the HDM annihilation channels contain it as one of the final dark states. It can then decay into the SM and account for most of the astrophysical signals. If such decays are predominantly to $\mu^+\mu^-, e^+e^-$ and sufficiently dominate over the dark photon decays from HDM annihilation, then the photon bounds from HESS can be satisfied. More detailed study is needed to determine the viable parameter space. An obviously safe parameter region is where $m_{\gamma_D} > 2m_0$, so that the dark photon is *not* dark-stable.

Let us now turn to this lightest particle of the dark spectrum, a scalar. We should verify that it decays efficiently enough in order to produce a signal for PAMELA and ATIC and that it decays fast enough to evade the BBN bounds. Here we will explicitly show that ϵ cannot do the job by itself and we indeed need other couplings between the dark and visible sectors, such as the effective operator (2.9) discussed in section 2.

To see this consider first the theory where ϵ alone connects the visible and dark sectors. The dark superpotential is given by

$$W_{\text{eff}} = \lambda T\bar{T}S + \frac{\kappa}{3!}S^3 + \eta TA\bar{X} + \bar{\eta}\bar{T}AX + \lambda T\bar{T}S + \mu_A A^2 + \mu X\bar{X} \tag{4.2}$$

The light fields T, $\bar{T}$, S by themselves couple with an accidental $\mathbb{Z}_3$ symmetry under which they transform as

$$\bar{T} \to e^{\frac{2\pi i}{3}}\bar{T}, \ S \to e^{-\frac{2\pi i}{3}}S, \ T \to T \ . \tag{4.3}$$

This symmetry survives T condensation and forbids the decay of the light scalar made from $S, \bar{T}$. The decay therefore requires an HDM loop since it is couplings to these heavy states that break the symmetry. The dominant decay process is then given by figure 4, of course when the external leg of T is replaced by its VEV. (One cannot directly decay to

Figure 4. An effective operator, responsible for the decay of $\bar{T}$ is formed at two-loop level (supergraph notations have been used). The insertions on the photonic lines denote ϵ and the cross denotes μ-term insertion.

SM massless photons since the ϵ vertex is proportional to $p^2_{photonline}$ which would vanish for on-shell photons.) A rough estimate yields a lifetime of order $\tau \sim 1000$ years.

A particle with such a long lifetime will decay into the charged SM particles long after BBN is over. Such decays would trigger EM showers in the plasma of the Early Universe and cause new processes of creating and destroying light elements. As has been shown in [32] the decays of our light scalar will cause severe overproduction of deuterium.

This problem can easily be circumvented when there are new, even very weak, couplings between visible and dark sectors, such as (2.9). For a large range, $\Lambda < 10^{11}$ GeV, assuming such couplings have coefficients which are roughly the same as the corresponding Yukawa couplings, a scalar $S \lesssim 600\,\mathrm{MeV}$ would decay to a muon pair in less than a second, before the start of BBN.[5] The muons in turn would decay to the electrons and positrons of the PAMELA/ATIC signals.

Of course, we must also estimate unwanted hadronic decays of S. Antiprotons are kinematically forbidden for such light S, but decays to π^0's would produce photons. Such decays would proceed via the coupling to strange quarks, $S\bar{s}s$, comparable to the $S\bar{\mu}\mu$ coupling if there is a rough proportionality to SM Yukawa couplings. At leading order in chiral perturbation theory $m_s\bar{s}s \equiv m_K^2|K|^2$, but kaon pairs are kinematically disallowed in S decays. Pairs of π^0 are not interpolated by $\bar{s}s$ at leading order in chiral perturbation theory but they do appear at the next order in the derivative expansion, which by naive dimensional analysis yields

$$m_s\bar{s}s \sim \frac{m_K^2}{(4\pi f_\pi)^2 f_\pi^2}(\partial_\nu \pi^0)^2. \tag{4.4}$$

(The non-derivative coupling of the pions is subdominant if S is heavier than the pion pair.) We then arrive at the branching ratio estimate

$$\frac{\Gamma_{S\to\pi^0\pi^0}}{\Gamma_{S\to\bar{\mu}\mu}} \sim \frac{1}{16}\left(\frac{M_{\mathscr{L}}}{M_{\mathscr{D}}}\right)^2\left(\frac{m_s}{m_\mu}\right)^2. \tag{4.5}$$

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[5]Analysing more carefully the constraints from BBN we notice that this demand can be even further relaxed. The authors of [32] impose strong bounds due to the late photo-dissociation processes only for particles with lifetime $\tau \gtrsim 10^4$ sec. Another constraint, which comes from entropy production considerations, does not apply in our case due to relatively small number density of the unstable relic. This means, that the scale Λ can be even as high as 10^{15} GeV.

The factor of 16 arises from the enhancement of the muon channel due to decay into distinct Dirac fermions rather than identical real scalars, and the derivative coupling of the scalars. The rough proportionality in the operators (2.10) to the SM Yukawa couplings leads to the appearance of the lepton to quark mass ratio in (4.5). The $M_{\mathscr{L}}/M_{\mathscr{D}}$ ratio accounts for the possibility of different suppression scales in equtation (2.9) for quarks and leptons.

We should bear in mind that this is just a rough estimate, depending on our assumption that the couplings (2.9) are proportional to SM Yukawa couplings. But in any case the free parameter $M_{\mathscr{L}}/M_{\mathscr{D}}$ can safely suppress this branching ratio. This provides one simple explanation for the excesses observed in PAMELA and ATIC, without attendant excesses of photons or antiprotons.

Of course, for scalar S mass below about $300\,\mathrm{MeV}$, the $\pi^0\pi^0$ is suppressed simply by phase space considerations, and for S mass below $200\,\mathrm{MeV}$ e^+e^- channel strongly dominates. These are also viable possibilities. But as one can see from figures 2, 3 they arise in rather small parts of our parameter space. It may however be more readily occurring in the part of the parameter space in which there are more dark VEVs than just $\langle T \rangle$, although we have not done this analysis.

Finally, we mention that the heavier dark fermions (R-odd) can also be kinematically dark-stable, and therefore can also decay to the SM. In particular, they can decay to the LDM and a lepton pair via a virtual exchange of the lightest dark scalar. Its lifetime is only an inverse loop factor larger than that of the lightest scalar, and hence poses no additional cosmological problems. It contributes to PAMELA/ATIC signals and therefore dilutes the proportion of these accounted for by dark photon decays.

4.3 Decays into the SM: collider signatures

Now let us turn to the possible collider signatures of this kind of model. Since our DM is not charged under the SM and is heavy, it is unlikely that any decay chain at the LHC will end up in the HDM. Thus the path at colliders to the dark sector is via the lightest MSSM R-odd particle, likely a neutralino.[6] This neutralino will cascade decay via the supersymmetric gauge kinetic mixing to the LDM particles, typically including dark-stable particles. If the dark photon is among these particles, it will promptly decay back to the SM, including the striking "lepton jet" signals predicted in [17].[7]

As discussed above, in our model the dark photon might also not be dark-stable. In that case the dark-stable particle decays back to the SM will be dominated by the lightest dark scalar. Note that the non-renormalizable coupling (2.9) that mediates such decays in our model is capable of a wide range of decay lengths. For different values of Λ, they are consistent with lepton jets, lepton jets with displaced vertices, and decay entirely outside the detector, while still being safe from BBN constraints.

The flavor structure of the couplings (2.9) is critical to understanding of which of these possibilities is more likely. The reason is that after electroweak symmetry breaking,

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[6]Since the lightest MSSM R-odd particle is unstable in this model, it can also be charged or colored. Decays of such a particle will therefore result in an extra charged or colored SM state beyond the possible lepton jet.

[7]See [33] for discussion of possibility measuring the "lepton-jets" in colliders

S exchange can mediate flavor-changing neutral current (FCNC) processes, which are very strongly constrained by experiments. If we assume only a rough proportionality to SM Yukawa couplings in (2.9), as we have mostly assumed earlier, then adequate suppression of FCNCs from S-exchange implies an S decay length of at least a meter. If this bound is saturated, then many decays of the lightest R-odd sparticle in the MSSM will end in lepton signals within the detector. Of course this is the edge of the viable parameter space, and the decay length could be considerably larger, in which case this source of lepton-jets is not realized within the detector.

If the S couplings in (2.9) are completely anarchic (comparable couplings for all pairs of SM flavors compatible with gauge invariance), then the S decays may still be exclusively to leptons if $m_S < 2m_\pi$, as required to avoid HESS bounds on photons from HDM annihilations. The couplings can be weak enough to satisfy FCNC bounds in hadrons as well as lepton-flavor violation constraints. Nevertheless, S's produced via MSSM sparticle decays can decay to striking μe pairs!

If for some reason there is precise proportionality to SM Yukawa couplings (Minimal Flavor Violation) in the S couplings of (2.9), then FCNC constraints are greatly weakened [34], in which case prompt scalar decays to leptons are robustly possible at colliders.

Apart from lepton-jet signals, the weak dark-visible couplings can also be probed indirectly by their contributions to rare low-energy processes [17, 35].

4.4 Dark matter abundance

Of course, any viable DM model must account for its measured cosmological abundance. Fitting this does not overconstrain our model, or alter the expected range of phenomenology, as we explain here. Note that the dark gauge coupling g_D dominates the annihilations of HDM. Thus with all other parameters of our model fixed, we can still fit the DM abundance by tuning g_D. We have not studied this in detail, but by the usual arguments in favor of WIMP DM, this tuned value will be $g_D \lesssim 1$.

We have presented our dark spectrum plots for only certain g_D values, but these spectra can be realized for whatever the fitted value of g_D is. Changes in g_D can be offset by the other parameters that feed into the scales of light dark matter, namely $m_{\frac{3}{2}}, \epsilon, \eta_i, \lambda, \kappa$. The masses and splittings in the HDM sector are independently controlled by c_i and μ_A (see (2.11), (2.12), (4.2)) respectively.

The DAMA signal, presumed to be accounted for by the iDM mechanism, depends on ϵ, m_{γ_D} for the scattering amplitude with ordinary matter, and the DM abundance and distribution for its magnitude. Since we are considering co-existing DM, where the iDM species is just one component, the relative abundance of this species can be varied by varying its mass (c_i-coefficient). This gives us one free parameter to fit the strength of the DAMA signal.

5 Conclusions and outlook

In this paper we introduced a framework which can accommodate the picture of the dark sector suggested in [16]. If the dark sector is sequestered from the SUSY-breaking hidden

sector, one can easily give an electroweak scale mass to the DM, using the generalized Giudice-Masiero mechanism, and the scale $\frac{m_{\frac{3}{2}}}{16\pi^2} \sim$ GeV naturally emerges for the dark gauge boson.

We took advantage of this framework to build an explicit model of the dark sector. We checked the viability of the model, verifying compatibility with the cosmological and astrophysical bounds. We emphasize that generic models of the dark sector contain additional stable particles and long-lived particles. Checking that the first is not overproduced while the second does not violate the predictions of BBN is necessary and can only be performed within concrete models.

Though our model has been built closely following the lines of the original proposal of [17], it possesses some distinct features. First, in our model the particle which is responsible for the Sommerfeld enhancement is not the particle whose decays dominate the PAMELA signal. We stress that the decaying particle can naturally be significantly lighter than the dark photon, successfully evading stringent bounds imposed by gamma-rays experiments. Second, the kinetic mixing between the visible and dark photons in our construction is not sufficient in order to produce a viable PAMELA signal, and additional couplings are identified. Such couplings may be a generic feature of such models. We also note that these additional couplings between sectors may or may not preserve the lepton jet signature at colliders, with little constraint from purely theoretical considerations.

We have focused here in demonstrating that our models are compatible with the various data, as we understand them currently. We have tried to do this using robust and plausible SUSY dynamics, in an easily generalizable and modular fashion. It would be interesting to conduct more precise numerical analysis of our models, and to precisely determine experimentally preferred regions in the parameter space, and to think through other possible experimental tests of models in this class. It would also be interesting to generalize our constructions, say to non-Abelian dark gauge group.

Acknowledgments

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