Non-Abelian Dark Sectors and Their Collider Signatures

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ABSTRACT: Motivated by the recent proliferation of observed astrophysical anomalies, Arkani-Hamed *et al.* have proposed a model in which dark matter is charged under a non-abelian "dark" gauge symmetry that is broken at ~ 1 GeV. In this paper, we present a survey of concrete models realizing such a scenario, followed by a largely model-independent study of collider phenomenology relevant to the Tevatron and the LHC. We address some model building issues that are easily surmounted to accommodate the astrophysics. While SUSY is not necessary, we argue that it is theoretically well-motivated because the GeV scale is automatically generated. Specifically, we propose a novel mechanism by which mixed D-terms in the dark sector induce either SUSY breaking or a super-Higgs mechanism precisely at a GeV. Furthermore, we elaborate on the original proposal of Arkani-Hamed *et al.* in which the dark matter acts as a messenger of gauge mediation to the dark sector. In our collider analysis we present cross-sections for dominant production channels and lifetime estimates for primary decay modes. We find that dark gauge bosons can be produced at the Tevatron and the LHC, either through a process analogous to prompt photon production or through a rare Z decay channel. Dark gauge bosons will decay back to the SM via "lepton jets" which typically contain > 2 and as many as 8 leptons, significantly improving their discovery potential. Since SUSY decays from the MSSM will eventually cascade down to these lepton jets, the discovery potential for direct electroweak-ino production may also be improved. Exploiting the unique kinematics, we find that it is possible to reconstruct the mass of the MSSM LSP. We also present several non-SUSY and SUSY decay channels that have displaced vertices and lead to multiple leptons with partially correlated impact parameters.

1. Introduction

Several intriguing observational results from high-energy astrophysics have motivated an exciting new proposal [1] in which a WIMP-like dark matter (DM) particle at 500-800 GeV annihilates primarily into leptons and is charged under a new "dark" force carrier. ATIC [2] detects an abundance of cosmic ray electrons between $300 - 800$ GeV, while PAMELA [3] sees an excess of positrons (but not anti-protons [4]) at 10-100 GeV. Together with the CMB haze [5, 6, 7], these observations paint a consistent picture whereby DM annihilates primarily into muons and/or electrons [8].

There are two sources of tension between these results and more conventional models of WIMP dark matter. First, assuming thermal freeze-out, the standard relic abundance calculation implies an annihilation cross-section that is at least a hundred times too small to explain the lepton excesses observed in astrophysical experiments. A "boost factor," typically attributed to local over-densities of dark matter, is often evoked in this case. A second difficulty is the non-observation of corresponding excesses in anti-protons [4] and gamma rays [9], which puts strong bounds on hadronic channels that are present in many dark matter models.

Motivated by the above considerations, the authors of Ref. [1] outline a scenario in which these apparent contradictions are reconciled. They introduce a 500-800 GeV WIMP that couples to a GeV scale dark gauge boson that kinetically mixes with the photon of the Standard Model (SM) [10][1] (see Ref. [12] for another recent suggestion with similar ingredients). A schematic illustration of this scenario is presented in Fig. 1. The ATIC and PAMELA data are explained by DM annihilation into the dark gauge boson which subsequently decays into electrons and muons. Elegantly enough, the $\mathcal{O}(1)$ GeV scale plays two independent roles. First, the new dark force carrier at $\lesssim$ GeV introduces a Sommerfeld enhancement [13, 14, 15, 16], giving a boost factor of the right size to enhance the DM annihilation cross-section[2]. Second, the absence of anti-protons in the PAMELA observations is now simply a result of kinematics [18].

The gauge group, G_{dark}, is a priori unspecified. However, it was observed in Ref. [1] that a non-abelian G_{dark} nicely accommodates the excited dark matter (XDM) [19] and inelastic dark matter (iDM) [20] mechanisms. XDM was proposed in order to explain the INTEGRAL [21] measurement of the 511 keV gamma-ray line at the center of the galaxy. The iDM scenario can accommodate the DAMA/LIBRA measurement of WIMP-nuclei scattering with other direct detection experiments [22, 23, 24]. Both XDM and iDM are non-standard WIMP scenarios in which a DM ground state can transition to and from new excited states via the emission of some field that couples back to the SM. If the DM lives in a multiplet of a non-abelian G_{dark}, then these ground and excited states can be the components of this multiplet, and transitions will emit dark gauge bosons that couple weakly to the SM. Independent of the

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[1]Ref. [11] analyzes particle physics bounds on such a light vector field and its possible connection to the HyperCP anomaly.

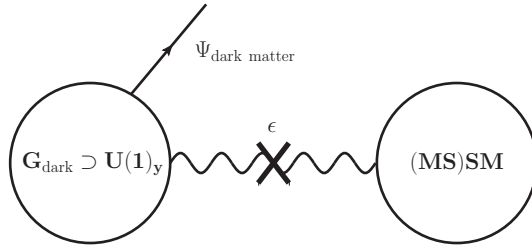[2]See Ref. [17] for an alternative, but related way for producing a large boost factor.

Figure 1: A schematic illustration of the minimal setup we consider in this paper. The dark sector and the SM are connected through kinetic mixing term suppressed by $\epsilon \lesssim 10^{-3}$. The dark matter multiplet may or may not couple directly to the SM. Supersymmetric extensions of this scenario are also discussed.

results from INTEGRAL and DAMA, we find the possibility of a non-abelian dark sector to be intriguing in its own right, with direct implications for the collider phenomenology. Thus throughout this paper we consider a dark sector with a non-abelian gauge symmetry that is completely broken by some dark Higgs sector[3].

In Section 2, we construct a catalog of explicit minimal models. Since G_{dark} needs to include a $U(1)$ factor for kinetic mixing with SM hypercharge, we take $G_{\mathrm{dark}} = SU(2) \times U(1)$. Our models differ only in their dark Higgs sectors, which are constructed to break G_{dark} completely and induce all the necessary couplings between the different states of the DM multiplet.

In Section 3, we discuss the mass splittings between the dark matter states. In order to obtain the small mass splittings needed for XDM and iDM, we consider DM that is a doublet or a triplet under $SU(2)_{\mathrm{dark}}$. The splittings may be generated radiatively from dark gauge boson loops. Another possibility is to generate them through higher-dimensional couplings between the dark matter and a single dark Higgs.

In Section 4, we consider the addition of SUSY to the dark sector. We observe that the minimal assumption of kinetic mixing between dark sector and SM hypercharge generates an effective FI term in the dark sector that is naturally of the desired scale, $\mathcal{O}(\mathrm{GeV})$. This term can break SUSY, or even more interestingly can generate a super-Higgs mechanism that leaves a supersymmetric dark sector with a ~ 1 GeV gap. Both of these scenarios typically result in light fermions that may have an influence on collider physics. We emphasize that this is a *leading* contribution which must be included in any SUSY scenario that includes kinetic mixing. Furthermore, within this scheme the DM can easily be a SM singlet, and so DM annihilations do not produce SM $W^{\pm}$ bosons that would dangerously decay to anti-protons that have not been observed by PAMELA. We also investigate the gauge mediation scenario originally proposed in Ref. [28] where DM is charged under the SM gauge group. An additional complication we address arises because SUSY restricts the form of the scalar

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[3]There are strong astrophysical constraints on a long range interaction from unbroken gauge symmetry with an unsuppressed coupling [25, 26, 27].

potential which is responsible for breaking G_{dark} completely. We provide several examples to overcome this difficulty.

In Section 5, we present several benchmark models for the dark Higgs sector. The resulting spectra of light vector bosons and scalars are explicitly computed and the relevant couplings are discussed.

In Section 6, we investigate the collider signatures of these models[4]. The kinetic mixing is the essential gateway to produce and observe dark sector states. Dark gauge bosons can be produced in processes analogous to prompt photon production in the SM. They can also be produced through rare Z decays. The dark sector states themselves dominantly decay into multiple $e^{\pm}$ and $\mu^{\pm}$, which are highly collimated and dubbed "lepton jets" [28]. Due to the non-abelian structure of G_{dark}, these lepton jets typically contain more than two leptons each. We discuss the observability of such signals at the LHC and at the Tevatron. We find that the cascade decays in the dark sector may result in displaced vertices or possible correlations between the $\slashed{E}_T$ and the lepton jet. Several such displaced vertices will produce uncorrelated impact parameters of decay products. In the case where the dark sector is supersymmetric, then it may be possible to detect direct electroweak gaugino production with enhanced reach both at the Tevatron and at the LHC. As a bonus, we find that we can exploit these cascades to perform absolute mass measurements of MSSM gauginos. Section 7 contains our conclusions.

Finally, let us briefly comment on the notational conventions used in this paper. In general, symbols referring to elements of the SM will be capitalized—so for example the SM hypercharge gauge coupling, gauge field, and field strength will be denoted by g_Y, B_μ and $B_{\mu\nu}$. In contrast, lowercase symbols will refer to elements of the dark sector, so the dark sector gauge coupling, gauge field, and field strength will be denoted by g_y, b_μ and $b_{\mu\nu}$. We will use h, or h' to denote dark Higgses. We denote the dark matter states by Ψ and we denote the SM and dark photon by γ and γ', respectively.

2. The Dark Sector and Symmetry Breaking

Let us begin by discussing the basic structure of the dark sector models that we will consider in this paper. We take the DM to be the lightest (and stable) component of some multiplet of the non-abelian group G_{dark}. As we will discuss in Section 3, such a multiplet is necessary if we wish to explain the INTEGRAL and/or DAMA signals along the lines of the XDM and iDM proposals of [19, 20, 29].

Furthermore, we follow the proposal of [1] in which the SM is coupled to the dark sector via a kinetic mixing term between SM hypercharge and a dark sector $U(1)$ gauge field:

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4}B_{\mu\nu}B^{\mu\nu} - \frac{1}{4}b_{\mu\nu}b^{\mu\nu} + \frac{\epsilon}{2}B_{\mu\nu}b^{\mu\nu} \tag{2.1}$$

where $B_{\mu\nu}$ and $b_{\mu\nu}$ are the SM and dark sector hypercharge field strengths, respectively. Because this marginal operator preserves all of the symmetries of the SM, it is relatively

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[4]We leave any precise matching to astrophysical observations for future work.

unconstrained phenomenologically. For a detailed analysis of kinetic mixing and the couplings it induces between SM and dark sector fields, see appendix A.

Since G_{dark} must contain a $U(1)$ factor[5], the minimal choice is of course $G_{\text{dark}} = SU(2) \times U(1)$. Furthermore, if G_{dark} is broken completely at a scale of $\sim$ GeV, then the resulting mass gap will relieve constraints from BBN on the number of relativistic degrees of freedom. However, in order to fully break charge, it is necessary to appropriately engineer a dark Higgs sector. As we shall see shortly, these scalars must also break a custodial $SU(2)$ in order to be phenomenologically viable. The necessity of breaking these symmetries demands a fairly elaborate dark Higgs sector.

First, let us consider the issue of charge breaking. Even for the simplest two Higgs doublet model, the criterion for charge breaking is quite complicated [30], for theories with more exotic Higgs representations, the space of charge breaking vacua is not even known. In appendix B, we present a straightforward method for deriving necessary conditions for charge breaking in two higgs doublet sectors, which we applied in order to obtain viable dark sector benchmark models.

Now let us explain the problem of the custodial symmetry. In the spirit of [1], we will assume that the DM is a multiplet of G_{dark} whose components are split in mass. The resulting excited and ground states have transitions mediated by dark gauge bosons that need to couple to the SM electric current if they are to realize the XDM and/or iDM scenarios (see Section 3). However, this mixing can be forbidden by a custodial symmetry of the dark Higgs sector. To see why this is the case, consider a model of two scalar doublets. We define $w_{a\mu}$ and b_{μ} to be the gauge bosons of G_{dark}, where b_{μ} is the abelian field which mixes with the SM hypercharge, B_{μ}. Assuming arbitrary vevs for the scalars, G_{dark} is broken, and in the $\{w_1, w_2, w_3, b\}$ basis, the gauge boson mass matrix takes the form

$$M^2_{\text{dark gauge}} = \begin{pmatrix} m_w^2 & 0 & 0 & \Delta_1 \\ 0 & m_w^2 & 0 & \Delta_2 \\ 0 & 0 & m_w^2 & \Delta_3 \\ \Delta_1 & \Delta_2 & \Delta_3 & m_b^2 \end{pmatrix} \tag{2.2}$$

As a consequence of the custodial symmetry present in any theory of only scalar doublets, the diagonal entries w_i are all equal. Applying a custodial $SU(2)$ transformation, we can rotate the components Δ_i completely into the w_3 direction. This yields a mass matrix which has a manifest $U(1)$ symmetry that acts as a phase rotation on $w_\pm = w_1 \pm iw_2$ (note that the gauged "electromagnetism" can still be broken while preserving this $U(1)$). Under this $U(1)$ the components of the DM multiplet have distinct charges—consequently the gauge bosons that mediate transitions among these states must also be charged, so they can only be the

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[5]It is actually possible to achieve mixing without an abelian factor through an S parameter type operator $\text{Tr}\,[\Phi w_{\mu\nu}]\,B^{\mu\nu}$, where Φ is some operator that transforms as an adjoint of the non-abelian group. In this paper we keep the abelian factor in order to investigate the collider signatures of the more general gauge group structure and ignore the existence of such operators. That is certainly justified in the case where no fundamental adjoints are present and the contribution is subleading.

$w_\pm$. However, $w_\pm$ have no components in the b direction, so they do not kinetically mix with SM hypercharge and thus cannot decay to SM particles.

Because the custodial symmetry is broken explicitly by the dark hypercharge, the couplings to the SM that are excluded at tree-level by this symmetry will be generated at one loop. Indeed, this may actually be desirable, since it generates an effective coupling for the DAMA transition that is suppressed beyond the ϵ^2 from the kinetic mixing. Another possibility, considered below, is to include additional Higgses that break custodial symmetry at tree-level.

In the case of SUSY models, we will be forced to significantly enlarge the Higgs sector. This is because many of the difficulties that arise in the non-SUSY case are exacerbated with the additional constraints imposed by SUSY. Moreover, in SUSY, all scalars are complex, which forces us to promote real Higgs triplets to complex Higgs triplets. This, along with the constraint of anomaly cancellation implies somewhat of a proliferation of Higgses in these theories.

In what follows, we enumerate several types of scalar sectors that break dark charge as well as custodial symmetry. We focus on models with the intention of later extending them with supersymmetry.

2.1 Doublet Models

A theory of one Higgs doublet is incapable of breaking charge, so we consider two doublets h_1 and h_2 with quantum numbers $\mathbf{2}_{-1/2}$ and $\mathbf{2}_{1/2}$ under $G_{\mathrm{dark}} = SU(2) \times U(1)$. A general renormalizable scalar potential that breaks charge is given by[6],

$$V(h_1, h_2) = \frac{\lambda_1}{2}\left(|h_1|^2 - |v_1|^2\right)^2 + \frac{\lambda_2}{2}\left(|h_2|^2 - |v_2|^2\right)^2 \tag{2.3}$$
$$+ \lambda_4 \left|h_1^T \epsilon h_2 - v_1 v_2 \cos\alpha\right|^2 + \lambda_3 \left(|h_1|^2 - |v_1|^2\right)\left(|h_2|^2 - |v_2|^2\right)$$

with,

$$\lambda_1 > 0, \quad \lambda_2 > 0, \quad \sqrt{\lambda_1 \lambda_2} + \lambda_3 > 0, \qquad \lambda_4 > 0 \tag{2.4}$$

v_2 is complex and charge is broken when $0 < |\cos\alpha| < 1$.

In the MSSM the conditions of Eq. (2.4) are violated at tree level. From the D-term contributions to the scalar potential we have $\lambda_1 = \lambda_2 = -\lambda_3 = (g^2 + g'^2)/8$. The inequality is saturated and the potential in Eq. (2.3) degenerates and contains a flat direction. To avoid such flat directions in the MSSM one must usually evoke a condition on the quadratic terms in the potential. Such potentials cannot be placed in the form of Eq. (2.3) and charge is not broken. Therefore the usual supersymmetric two doublets model will not suffice and we need additional contributions to the scalar potential in order to satisfy the condition, Eq. (2.4).

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[6]This is not the most general renormalizable scalar potential possible. One can add three more terms, $|h_1|^2|h_2|^2$, $|h_1|^2 h_1^T \epsilon h_2$, and $|h_2|^2 h_1^T \epsilon h_2$ which are consistent with all the symmetries. However, these simply complicate the potential and are not required for breaking charge. A more general analysis of the vacuum structure can be found in Ref. [30]

In addition, since this model has a custodial symmetry, it fails to have proper mixings between the gauge bosons. Nonetheless, since the custodial symmetry is broken by dark hypercharge, the gauge boson mixing receives one-loop radiative corrections that break the custodial symmetry. From this point of view there is also no reason not to include higher dimension custodial violating operators that can be generated if heavy (triplet) states have been integrated out. In fact, we include such irrelevant operators in the benchmark model of Section 5.1.1.

Here we also note the presence of an unfortunate $\mathbb{Z}_2$ symmetry that is present in the $\tan\beta = 1$ limit. This symmetry needs to be broken since it prevents two of the dark gauge bosons from coupling to SM electric charge (see appendix C).

2.2 Doublet/Triplet Models

An obvious way to break custodial $SU(2)$ at tree-level is to augment the two doublet model with a light triplet of $SU(2)$. For instance, consider a model of one doublet, h, and one real triplet, Φ, with dark quantum numbers $\mathbf{2}_{1/2}$ and $\mathbf{3}_0$, respectively. In order to realize a charge breaking angle between the doublet and triplet, we include the following two operators: $h^\dagger \Phi h$ and $h^T \epsilon \Phi h$. Since the latter has nonzero hypercharge, we must multiply it by a new hypercharged singlet, S, in order to include it in the potential:

$$V(h, \Phi, S) = \frac{\lambda_h}{2} \left(|h|^2 - |v_h|^2 \right)^2 + \frac{\lambda_\Phi}{2} \left(\mathrm{Tr}\left[\Phi\Phi \right] - |v_\Phi|^2 \right)^2 + \frac{\lambda_S}{2} \left(|S|^2 - |v_S|^2 \right)^2 \quad (2.5)$$
$$+ \, c_1 h^\dagger \Phi h + (c_2 S h^T \epsilon \Phi h + \text{h.c.})$$

Alternatively, we might consider a model with two doublets and one triplet. This is more natural if we wish to eventually include SUSY. The scalar potential takes the form:

$$V(h_1, h_2, \Phi) = V(h_1, h_2) + \frac{\lambda_\Phi}{2} \left(\mathrm{Tr}\left[\Phi\Phi \right] - |v_\Phi|^2 \right)^2 \quad (2.6)$$
$$+ \, c_1 h_1^\dagger \Phi h_1 + c_2 h_2^\dagger \Phi h_2 + \left(c_3 h_1^T \epsilon \Phi h_2 + \text{h.c.} \right)$$

where $V(h_1, h_2)$ is the contribution from doublets alone defined in eq 2.3.

We can impose an additional $\mathbb{Z}_2$ symmetry $\Phi \to -\Phi$, that forbids tree-level couplings between the triplet and doublets: $c_1 = c_2 = c_3 = 0$. This enhanced global symmetry implies the existence of two pseudo-Goldstone bosons which obtain masses at one-loop ~ 10 MeV. These pseudo-Goldstone bosons will be produced at the bottom of dark sector cascades. They decay into leptons through either two off-shell dark gauge bosons or at one-loop (see Fig. 10 and the discussion in Section 6.1.1). Either way, the long lifetime causes the pseudo-Goldstone boson to escape the detector at colliders. Since those are pseudo-scalars they will not contribute to the Sommerfeld enhancement of DM annihilations in the early universe and their mass is therefore not bounded by the limits derived in Ref. [31].

3. Dark Matter Mass Splitting

The authors of [1] observed that a DM multiplet of some non-abelian G_{dark}, given appropriate mass splittings, can in principle realize the XDM explanation of INTEGRAL [19] and also the iDM mechanism for reconciling the DAMA annual modulation with the null result of other direct detection experiments [20, 29]. In this section we briefly review these proposals, and discuss concrete ways of generating the appropriate mass splittings within concrete theories.

The INTEGRAL collaboration has provided an extremely refined measurement of the 511 keV line of positronium annihilation coming from the galactic center. In the XDM scenario, WIMPs in the galactic center scatter into an excited state, lying ~ 1 MeV above the ground state. The excited state then de-excites into e^+e^- which provides the excess positrons needed. In terms of model-building we need a splitting of ~ 1 MeV between two states in the DM multiplet. Transitions between these two states are mediated by a dark gauge boson with some component of the dark hypercharge (which in turn couples to SM leptons).

In contrast, DAMA is a direct detection experiment which seeks to measure the scattering of galactic WIMPS off of NaI(Tl). Assuming a standard WIMP with elastic scattering, several other experiment such as CDMS[22, 23], XENON [24], and ZEPLIN [32] exclude DAMA's measured annual modulation by many orders of magnitude. The iDM proposal reconciles these experiments by proposing that the WIMP can only scatter off of nuclei through an inelastic process by which the the DM is converted into a slightly excited state. Since the WIMP kinetic energy is fixed and the threshold for the inelastic transition is dependent on the atomic number of the nuclei, this iDM scenario can simultaneously predict a null result at CDMS and a positive result at DAMA[7]. Considering fermionic DM, this scenario can be accommodated by including a mass splitting of around ~ 100-150 keV [20, 29] between the lightest two Majorana states of the fermion. The bottom line for model building is that to evade CDMS and CRESST[33] bounds, the DM must be split from the next heaviest Majorana state by at at least 100 keV.

Before we consider mechanisms for generating the required splittings, we must ascertain that there is no elastic scattering which would have been seen in direct detection experiments. One possibility is to begin with Majorana dark matter in a real representation of the dark gauge symmetry. Gauge bosons then couple different states of the multiplet and radiative corrections, to be discussed in Section 3.1, can split the masses of these states. But if the dark matter begins in a complex representation, for example if it has dark or SM U(1) charge, then it must be Dirac-like at high-energies. Then the model is already excluded by direct detection experiments since the elastic scattering of Dirac-like dark matter is not sufficiently suppressed unless $\epsilon \lesssim 10^{-6}$. However, it is possible to split the masses of the Majorana components of the Dirac fermions by using the same scalar sector that is responsible for breaking dark gauge symmetry. For instance, if we imagine that ϕ is some scalar singlet whose vev breaks global

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[7]XENON and ZEPLIN, which both use Xe as a target should be able to exclude the iDM scenario, but at the moment these experiments are background limited [29]

fermion number and $U(1)_y$, we can add a term such as,

$$\mathcal{L}_{\text{Majorana}} = \phi \Psi \Psi + \phi^* \Psi^c \Psi^c \tag{3.1}$$

where Ψ and Ψ^c are the Weyl components of some DM multiplet. If ϕ develops a vev of order $\sim$ GeV, it will generate a Majorana mass splitting that forbids any elastic scattering and evades direct detection bounds.

Another possibility is to use a higher dimensional operator with a dark sector doublet, h,

$$\mathcal{L}_{\text{Majorana}} = \frac{1}{M_X} h \Psi \Psi h, \tag{3.2}$$

where $M_X \sim$ TeV. In this case, the Majorana splitting is of order $\sim$ MeV which again kinematically forbids elastic scattering. If DM is charged under the SM as part of a $\mathbf{5} + \bar{\mathbf{5}}$ multiplet, then a dimension 6 operator is required to contract both dark hypercharge and SM quantum numbers. For example, we can use the operator $1/M_X^2 \, H\Psi\Psi H\phi$, where H is the SM Higgs, ϕ is a singlet that soaks up Ψ's dark hypercharge and gets a vev at ~ 1 GeV, and $M_X \sim$ TeV. This leads to Majorana splitting of order $\sim$ GeV, which forbids elastic scattering.

Any of the possibilities mentioned above can be employed to evade direct detection from CDMS. In the next subsection, we consider two possible way for generating the appropriate ~ 1 MeV and ~ 100 keV mass splittings necessary for XDM and iDM.

3.1 Radiative Splitting

As is well-known [34], spontaneous symmetry breaking of a non-abelian gauge group generates radiative mass splittings within a multiplet of the symmetry. We take the DM multiplet to have mass $\sim 500 - 800$ GeV, and to be charged under the dark $SU(2) \times U(1)$. As discussed in section 2, realistic dark sectors must break charge and custodial $SU(2)$, but to develop some intuition about the radiative mass splittings, we will begin by considering the limit where these symmetries are preserved. In this limit the mass splittings among the multiplet take a particularly simple form,

$$\begin{aligned}
\Delta m_{ij} = & \frac{\alpha^{\text{dark}}}{2} (q_i^2 - q_j^2) M_z \\
& - \frac{\alpha_2^{\text{dark}}}{2} \left((T_i^3)^2 - (T_j^3)^2 \right) (M_z - M_w),
\end{aligned} \tag{3.3}$$

where we define α_2^{dark}, and α^{dark} as usual with respect to $SU(2) \times U(1)$ couplings. The charges are $q_i = T_i^3 + Y$ and T_i^3 is the ith eigenvalue of the third $SU(2)$ generator. In the more general limit where charge and custodial symmetry are broken, one must use the appropriate vector boson mass eigenstates and their couplings to the fermions in order to compute the mass correction (Eqs. D-1). This is a straightforward computation, however, in general it does not yield a simple analytic result. Nevertheless, it is clear that there are two factors which control
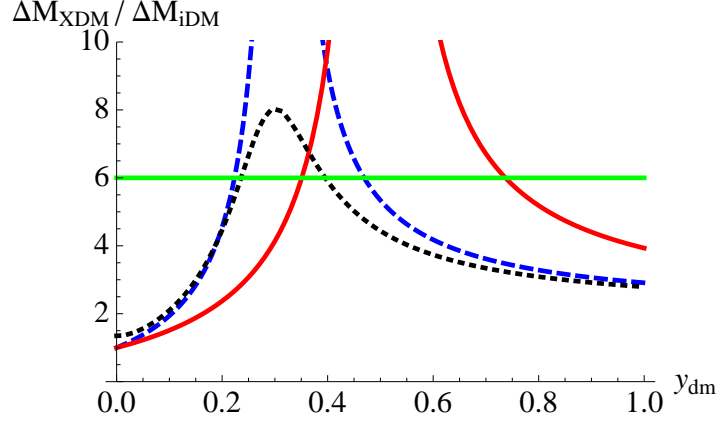
Figure 2: The ratio of the XDM splitting to the iDM splitting as a function of triplet dark matter $U(1)_y$ hypercharge. The green horizontal line indicates the minimum ratio for simultaneously achieving both splittings. Red (line) is an example of two Higgs doublets with charge preserved, blue (dashed) represents two Higgs doublets with charge broken, and black (dots) adds a Higgs triplet to the previous case. For this example, the gauge couplings are $g = 0.97$ and $g_y = 0.26$, and in terms of Eq. 2.3 we have for all three models $v_1 = 0.9$ GeV, $v_2 = 1.1$ GeV and $\lambda_{1,2,3,4} = 1$. Red (line) and black (dots) add charge breaking with $\cos\alpha = 0.75$, and for black (dots), in terms of Eq. 2.6, $\lambda_\Phi = 1$, $v_\Phi = 1$ GeV and the triplet is decoupled from the doublets at tree-level by imposing the discrete symmetry: $\Phi \to -\Phi$.

the mass splitting: first, the differences between masses of the vector bosons; and second, the couplings of the different members of the representation to the vector bosons.

As a simple example with all the required splittings and couplings we can consider a triplet with hypercharge $y = 1/2 - \delta$. We generate both large, $\Delta M \sim \alpha^{\text{dark}} M_z$ and small $\Delta m \sim \delta \alpha^{\text{dark}} M_z$ splittings. The correct couplings to account for the XDM and iDM scenarios are induced when charge and custodial breaking corrections are included. A realistic model certainly need not be based on such odd charge assignments, however, this example serves to illustrate how straightforward it is to obtain the correct splittings and couplings. In Figs. 2 and 3, we consider some of the more general models of Section 2, which include charge and custodial symmetry breaking, and we plot the exact ratio of the two splittings relevant to XDM and iDM as a function of the parameters. The corrections induced in supersymmetric models are discussed in appendix D.

3.2 Mass splitting from higher dimensional operator

It is also possible to generate the INTEGRAL and DAMA mass splittings from higher dimension operators alone. The key observation is that $\delta m \sim \Lambda^2_{\text{dark}}/M_X \sim$ MeV, which is of the desired range.

As an example, we consider two Weyl fermions Ψ, Ψ^c which are $\mathbf{2}_{1/2}$ and $\bar{\mathbf{2}}_{-1/2}$ under G_{dark}. It is possible to achieve all the required splittings and transitions with a single scalar
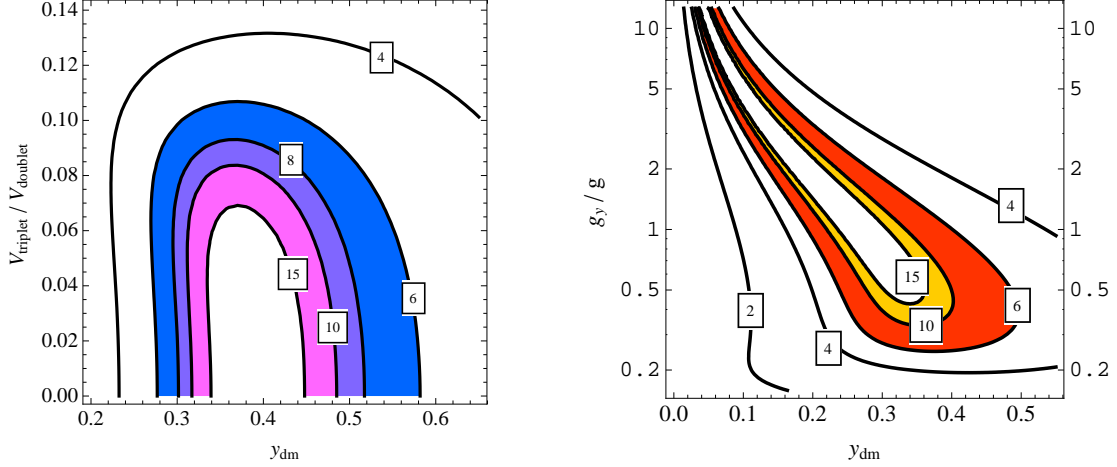
Figure 3: Two contour plots of the ratio of the XDM splitting to the iDM splitting for triplet dark matter with two Higgs doublets and one Higgs triplet. The shaded regions represent splitting ratios where XDM and iDM can be achieved simultaneously. In both plots, the horizontal axis is the dark matter $U(1)_y$ hypercharge. The vertical axis of the left plot represents the ratio of the triplet to doublet VEVs, v_Φ/v, where $v^2 = v_u^2 + v_d^2$ and $\langle\Phi\rangle = v_\Phi T_3$. The vertical axis of the right plot represents the ratio of dark hypercharge and $SU(2)$ couplings, g'/g. For both plots, the triplet is decoupled from the doublets at tree-level by imposing the discrete symmetry: $\Phi \to -\Phi$, and in terms of Eq. 2.3 we have $v_1 = 0.9$ GeV, $v_2 = 1.1$ GeV, $\cos\alpha = 0.9$ and $\lambda_{1,2,3,4} = 1$. For the left plot, the gauge couplings are $g = 0.97$ and $g_y = 0.26$. For the right plot, we have also chosen, in terms of Eq. 2.6, $\lambda_\Phi = 1$ and $v_\Phi = 1$ GeV.

doublet,

$$\mathcal{L} \supset M_\Psi \Psi \Psi^c + \frac{\lambda_1}{M_X}\Psi h \Psi h + \frac{\lambda_2}{M_X}\Psi^c h^c \Psi^c h^c + \frac{\lambda_3}{M_X}\Psi^c h^c \Psi h + h.c. \tag{3.4}$$

with $\Psi = (\psi_\nu, \psi_e)$ and $\Psi h \equiv \psi_i \epsilon_{ij} h_j$ and $h_i^c = \epsilon_{ij} h_j^*$. Once the scalar doublet gets a vev, $\langle h \rangle = (0, v)$, the "neutrino" components of Ψ and Ψ^c mix through the following matrix,

$$M = \begin{pmatrix} \lambda_1 \bar{v} & M_\Psi + \lambda_3 \bar{v} \\ M_\Psi + \lambda_3 \bar{v} & \lambda_2 \bar{v} \end{pmatrix}, \tag{3.5}$$

where $\bar{v} = v^2/M_X \sim$ MeV. In the limit where $\lambda_1 = \lambda_2 = 0$ the states are maximally mixed, $\psi^\pm = (\psi_\nu \pm \psi_\nu^c)/\sqrt{2}$, and form a Dirac pair of mass $M_\Psi + \lambda_3 \bar{v}$ which provides the XDM splitting when compared with the ψ_e, ψ_e^c states of mass M_Ψ.

With non-zero λ_1 and λ_2 we have,

$$\psi_1' = \cos\theta\psi^+ + \sin\theta\psi^- \qquad m_1 = M_\Psi + \frac{\bar{v}}{2}(2\lambda_3 + \lambda_1 + \lambda_2) \tag{3.6}$$

$$\psi_2' = -\sin\theta\psi^+ + \cos\theta\psi^- \qquad m_2 = -M_\Psi - \frac{\bar{v}}{2}(2\lambda_3 - \lambda_1 - \lambda_2)$$

with $\sin\theta \approx (\lambda_1 - \lambda_2)\bar{v}/4M_\Psi$. The mass difference between the two states is $|\Delta m_{12}| = (\lambda_1 + \lambda_2)\bar{v}$. So, by tuning λ_1 against λ_2 we can achieve $\Delta m_{12} \sim 0.1$ MeV $= 100$ keV as
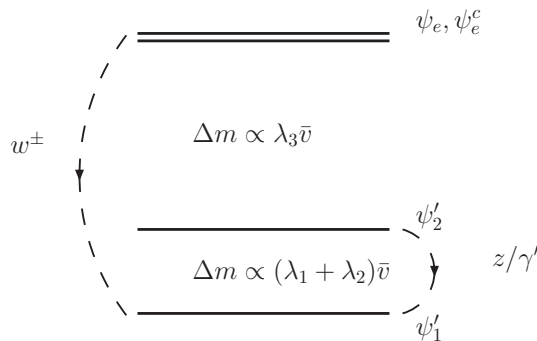
Figure 4: The resulting spectrum for a Dirac doublet with majoron coupling.

required by iDM. The coupling of the mass eigenstates to the dark gauge boson is given by,

$$
\begin{aligned}
g_y \bar{\Psi} \slashed{z} \Psi - g_y \bar{\Psi}^c \slashed{z} \Psi^c = {}& g_y \sin\theta \cos\theta \, (\bar{\psi}'_2 \slashed{z} \psi'_2 - \bar{\psi}'_1 \slashed{z} \psi'_1) \\
& - g_y \cos^2\theta \, \bar{\psi}'_1 \slashed{z} \psi'_2 + h.c.
\end{aligned}
\tag{3.7}
$$

In this case the ratio of the elastic to inelastic coupling is approximately $\sin\theta = (\lambda_1 - \lambda_2)(\bar{v}/M_\Psi) \sim 10^{-7}$, which is sufficiently suppressed. The spectrum relevant for this case is shown in Fig. 4.

4. Generation of the Dark Sector Mass Scale

As noted in [28], a particularly nice feature of a SUSY dark sector is that the GeV scale is naturally generated by gauge mediated SUSY breaking from the SM. In this section, we elaborate on this scenario in detail. Furthermore, we propose an even more minimal alternative in which "kinetic mixing mediation" breaks SUSY or induces a super-Higgs mechanism at a scale of several GeV in the dark sector. As we will discuss, these theories typically have light fermions which affect the collider physics.

For gauge mediation, dark matter itself can act as the messenger if we take it to be charged under the SM as part of a $\mathbf{5} + \bar{\mathbf{5}}$ multiplet. Dark matter annihilations then also produce SM electroweak gauge bosons, resulting in hadronic channels. But since the GeV scale can be generated by kinetic mixing mediation alone, there is no need to charge dark matter under the SM.

Although we focus on kinetic mixing mediation and gauge mediation for the rest of this section, and when we construct benchmarks in Section 5, there are other ways to break SUSY in the dark sector. We would like to stress that the rest of our paper, in particular the model-independent discussion of collider signatures in Section 6, does not depend on how SUSY is broken in the dark sector. One alternative is that there is high-scale gauge mediation and a GeV scale gravitino [28]. Then SUSY is broken in the dark sector at the GeV scale by a "Planck slop." Another possibility is that the dark matter mass is related to

the mechanism that sets the MSSM μ parameter, for example due to a superpotential of the form: $\lambda S H_u H_d + \lambda' S \Psi \Psi^c$. A vev for F_S is communicated to the dark sector through gauge mediation with dark matter as messengers.

4.1 SUSY Breaking from Kinetic Mixing

In [35] it was observed that mixing of gauge boson kinetic terms will induce mixed D-term contributions to the action that can communicate SUSY breaking between two sectors that are otherwise decoupled.[8] There the authors noted new, possibly dangerous contributions to SUSY breaking to the MSSM from this effect. In this section, we use this effect to our advantage in order to mediate SUSY breaking from the SM to the dark sector. We should note that while we can choose to make this the dominant mechanism for breaking SUSY in the dark sector, it is always present at the GeV scale.[9]

As was originally proposed in [1], we have been assuming that the dark sector and the SM are coupled via a marginal gauge kinetic mixing between the dark hypercharge, $U(1)_y$ and the SM hypercharge, $U(1)_Y$. If both $U(1)$'s are fundamental, then the kinetic mixing is a UV boundary condition, sensitive to physics at the highest scales. Instead, if either $U(1)$ is embedded in a GUT, then the kinetic mixing is only induced below the scale of GUT breaking by integrating out fields charged under both $U(1)$'s. In this case we can estimate its size. In particular, heavy fields charged under both the SM and the dark sector will induce a gauge kinetic mixing:

$$\mathcal{L}_{\text{gauge}} = \frac{1}{4} \int d^2\theta \left(W_Y W_Y + W_y W_y - 2\epsilon W_Y W_y + \text{h.c.} \right) \tag{4.1}$$

$$\epsilon \sim -\frac{g_Y g_y}{16\pi^2} \log\left(\frac{M^2}{M'^2}\right) \tag{4.2}$$

where g_y and g_Y are the gauge couplings for the dark and SM hypercharges, respectively, M and M' are the masses of components of the heavy particle multiplet. Assuming that these mass scales are not too separated and that the gauge couplings are of reasonable size, this gives an estimate of $\epsilon \sim 10^{-3} - 10^{-4}$. Interestingly, this not only gives the right scale to explain the DAMA cross-section, but also generates a scale of around a GeV in the dark sector. Eq. (4.1), along with the Kahler potential, implies a D-term potential:

$$V_{\text{gauge}} = \frac{1}{2}D_Y^2 + \frac{1}{2}D_y^2 - \epsilon D_Y D_y + g_Y D_Y \sum_i Q_i \left|H_i\right|^2 + g_y D_y \sum_i q_i \left|h_i\right|^2 \tag{4.3}$$

where H_i and h_i denote the SM and dark sector Higgs, respectively, and Q_i and q_i denote their SM and dark sector hypercharges. Integrating out the SM fields, H_i and D_Y, generates a cross term $\epsilon D_y \langle D_Y \rangle$ in the low-energy theory. Thus, in the infrared, this induces an effective

---

[8]Strictly speaking, this is a form of gauge mediation according to the definition of Ref. [36]

[9]Kinetic mixing mediation is neglected in some recent $U(1)$ dark sector papers, for instance Ref. [37] focuses on a form of mediation that is sub-leading in ϵ.

Fayet-Iliopolous D-term for D_y, for which

$$V_{\text{gauge}} \supset \epsilon D_y \langle D_Y \rangle = \xi D_y \tag{4.4}$$

$$\xi = \epsilon \langle D_Y \rangle = \epsilon \frac{g_Y}{2} \cos 2\beta \, v^2 \tag{4.5}$$

where in the last equality we have substituted in for $\langle D_Y \rangle$ from the MSSM. For $\epsilon \sim 10^{-3} - 10^{-4}$, ξ is at the GeV^2 scale. Thus, given the minimal assumption of kinetic mixing and SUSY, we obtain precisely the right scale to account for PAMELA and ATIC with the Sommerfeld enhancement.

With an effective FI term at low energies, it is straightforward to break SUSY in the dark sector. In particular, a generic superpotential for the dark Higgses, h_i, will break SUSY because the F and D terms cannot be simultaneously set to zero. While this SUSY breaking generates scalar soft masses, it does not generate soft masses for gauginos. Moreover, since SUSY is broken within the dark sector this will typically introduce a massless dark sector Goldstino[10], assuming the absence of explicit SUSY breaking operators. So, within this mechanism there are light fermions. We present a concrete model of this type in Section 5, and mention the possibility of associated missing energy signals in Section 6.1.2.

In the opposite extreme, we can take the superpotential to be less generic, or perhaps even trivial, and so the dark Higgs potential is dominated by D-terms. Here the dark Higgses simply align to set the dark hypercharge D-term to zero. In this limit SUSY is actually preserved in the dark sector, but a super-Higgs mechanism will generate a GeV scale dark sector that may still be consistent with a Sommerfeld enhancement and αm_z mass splittings for DM. A more minimal superpotential also can imply the existence of light pseudo-Goldstone fields and their superpartners. Finally, we note that unlike the SUSY breaking case, this super-Higgs scenario will also generate GeV scale gaugino masses.

4.2 SUSY Breaking From $5 + \bar{5}$ Messengers

In this section, we elaborate on the gauge mediation proposal of [28] in which a multiplet of $5 + \bar{5}$ messengers is charged directly under both the dark sector and SM, thereby communicating SM SUSY breaking to the dark sector. We consider the additional possibility that the lightest component of the messenger supermultiplet is in fact the DM.

Let us determine the various contributions which set the scale of masses for the scalar and fermion components of the DM $5 + \bar{5}$ multiplet. First, we assume that the fermions have a SUSY mass, $m_f^{(2,3)}$, that splits the doublet and triplet components. Second, in the case where low-scale gauge mediation explains SUSY breaking in the MSSM, then the doublet and triplet scalars of the $5 + \bar{5}$ receive identical soft mass contributions to that of the sleptons and right-handed down squarks of the MSSM. We denote this contribution by $m_s^{(2,3)} \sim 100$ GeV. Finally, the scalar DM can in principle also receive soft mass contributions from whatever dynamics set its μ and $B\mu$ terms, which we denote by $B\mu^{(2,3)}$. Instead of specifying these

---

[10]The gravitino will eat a linear combination of this field and the Goldstino associated with SUSY breaking in the MSSM

dynamics, we will choose a model-independent parameterization for the DM supermultiplet masses. The scalar doublets and triplets of the $\mathbf{5} + \bar{\mathbf{5}}$ have a scalar mass matrix given by:

$$M_{\mathbf{2,3}}^2 = \begin{pmatrix} \left[m_f^{(\mathbf{2,3})}\right]^2 + \left[m_s^{(\mathbf{2,3})}\right]^2 & B\mu^{(\mathbf{2,3})} \\ B\mu^{(\mathbf{2,3})} & \left[m_f^{(\mathbf{2,3})}\right]^2 + \left[m_s^{(\mathbf{2,3})}\right]^2 \end{pmatrix} \tag{4.6}$$

whose eigenvalues, $m_\pm^{(\mathbf{2,3})}$, are given by

$$\left(m_\pm^{(\mathbf{2,3})}\right)^2 = \left(m_f^{(\mathbf{2,3})}\right)^2 + \left(m_s^{(\mathbf{2,3})}\right)^2 \pm B\mu^{(\mathbf{2,3})}. \tag{4.7}$$

The lightest component of the doublet supermultiplet corresponds to dark matter, and we choose it to have mass $\sim 500 - 800$ GeV, which is favored by ATIC.

Note that the messenger supertrace of the $\mathbf{5} + \bar{\mathbf{5}}$ mass matrix is non-zero, and proportional to $m_s^{(\mathbf{2,3})}$. As discussed in [38], this non-zero supertrace generates a logarithmically UV sensitive soft mass for the dark sector scalars. Indeed, since the messenger supertrace is positive, this implies a negative soft mass for the dark Higgs:

$$m_h^2 \approx -8 \left(\frac{\alpha}{4\pi}\right)^2 \left(2\left[M_s^{(\mathbf{2})}\right]^2 + 3\left[M_s^{(\mathbf{3})}\right]^2\right) \log\left(\frac{\Lambda_{UV}^2}{m_f^2}\right) C_a S_q \tag{4.8}$$

where C_a is the dark scalar's quadratic Casimir, S_Q is dark matter's Dynkin index, and Λ_{UV} is set by the messenger scale of SUSY breaking to the SM. The negative soft mass squared allows for G_{dark} to break. This gives us a way to break the symmetries, independent of the effect of RGE running. It is our assumption that the contributions due to running are suppressed. For low-scale gauge mediation, $\Lambda_{UV} \sim 30 - 100$ TeV, and because of this logarithmic enhancement and the combined effect of five messengers, we find that our desired scale of $m_h^2 \sim 1$ GeV^2 implies that $m_s^{(\mathbf{2,3})} \sim 50$ GeV. This indicates a bit of tension numerically because we expect that $m_s^{(\mathbf{2,3})}$ is set by the SM soft mass scale of hundreds of GeV.

Additionally, if we want fermionic DM, then there is the additional constraint that the fermion is the lightest component of the dark matter supermultiplet: thus $(B\mu^{(\mathbf{2})})^{1/2} < m_s^{(\mathbf{2})} \ll m_f^{(\mathbf{2})}$. Since the DM $B\mu$ contribution breaks the dark sector R-symmetry, the gaugino soft masses are suppressed if we assume that the triplet component satisfies the same condition:

$$m_\lambda \approx \frac{\alpha}{2} S_q \left(2\frac{B\mu^{(\mathbf{2})}}{m_f^{(\mathbf{2})}} + 3\frac{B\mu^{(\mathbf{3})}}{m_f^{(\mathbf{3})}}\right) \tag{4.9}$$

This implies light gauginos and the generic prediction is that fermionic dark matter implies that the lightest dark sector particle is a mostly-gaugino fermion. This conclusion can be avoided by raising $B\mu^{(\mathbf{3})}$ while maintaining $(B\mu^{(\mathbf{2})})^{1/2} \ll m_f^{(\mathbf{2})}$.

The dark sector Higgses require GeV scale μ and $B\mu$ terms to help break dark gauge symmetry and lift runaway directions. These terms can be generated by additional dynamics

that communicate SM SUSY breaking to the dark Higgses, in general also resulting in new two-loop contributions to the dark scalar masses. We will assume that these contributions to m_h^2 are subdominant to the usual gauge mediation contributions of Eq. 4.8. A recent paper identifies a class of general gauge mediation models that satisfy this assumption [39].

Let us note that while Eqs. 4.8 and 4.9 are approximations, our benchmark model of gauge mediation in Section 5.2.2 employs the full expressions of Ref. [38].

5. Benchmark Models

In this section, we present four detailed benchmark dark sector models and their spectra. The models break dark gauge symmetry and custodial symmetry, generating the dark matter splittings and couplings necessary to explain the astrophysical data, as explained in Section 2. These examples illustrate some of the theoretical issues discussed above, and their spectra and couplings serve as starting points for thinking about the types of cascades that can occur in GeV scale dark sectors. We begin in Section 5.1 with two non-SUSY models, where the GeV scale is put into the scalar potential by hand. We then consider two SUSY examples in Section 5.2, where the GeV scale is generated radiatively in the dark sector from interactions with the Standard Model.

For each example we consider an $SU(2) \times U(1)_y$ dark sector and triplet dark matter. We take the Majorana components of the dark matter fermions to be split by enough to avoid direct detection bounds, for example by one of the mechanisms discussed in Section 3. We then calculate the radiative splittings among the triplet, induced by dark symmetry breaking, as in Section 3.1. We take the ground state to correspond to dark matter, the heaviest excited state to allow for the XDM explanation of INTEGRAL, and the first excited state to allow for the iDM explanation of DAMA. We allow complex parameters to carry imaginary parts in order to avoid unbroken CP symmetry in the dark sector which may lead to stable states. This is not necessarily a problem and may actually have additional interesting signatures, but we'd like to keep the spectrum as general as possible for the present discussion.

5.1 Non-SUSY Benchmark Models

5.1.1 Non-SUSY 1: Two Doublets

We begin with the two doublet model of Section 2.1, where h_1 and h_2 have dark quantum numbers $\mathbf{2}_{-1/2}$, $\mathbf{2}_{1/2}$. We have chosen a benchmark which breaks charge and radiatively generates the XDM and iDM splittings, and we have calculated its mass spectrum (Fig. 5). As discussed in section 2, the custodial $SU(2)$ symmetry of the Higgs sector determines the tree-level gauge boson spectrum. The gauge bosons that couple between the different dark matter mass eigenstates, $w_\pm$, do not mix with the b and are degenerate in mass. Custodial symmetry is broken at one-loop and in general due to higher-dimensional operators. The DAMA inelastic scattering is therefore suppressed relative to models where custodial symmetry is broken at tree-level. For this benchmark, we induce the iDM coupling by including

the dimension 6 custodial-breaking operators $c_1^T |h_1 D h_1|^2$ and $c_2^T |h_2 D h_2|^2$, with coefficients c_1^T and c_2^T expressing the loop-suppression.

For the benchmark, we choose the gauge couplings: $g = 0.46$ and $g_y = 0.19$. The dark matter hypercharge is chosen to be $y_{\text{dm}} = 1/2$. In the limit of small charge breaking, this choice leads to one small and one large dark matter splitting, as discussed in Section 3.1. In terms of the potential of Eq. 2.3, the parameters are: $v_1 = 1.5$ GeV, $v_2 = (1.5 + 3.2 i)$ GeV, $\lambda_1 = 0.5$, $\lambda_2 = 0.3$, $\lambda_3 = -0.031$, $\lambda_4 = 0.5$, and $\cos \alpha = 0.8$. The coefficients of the custodial-breaking dimension 6 operators are chosen to be $c_1^T = 2.8 \times 10^{-4}$ GeV^{-2} and $c_2^T = -5.7 \times 10^{-4}$ GeV^{-2}.
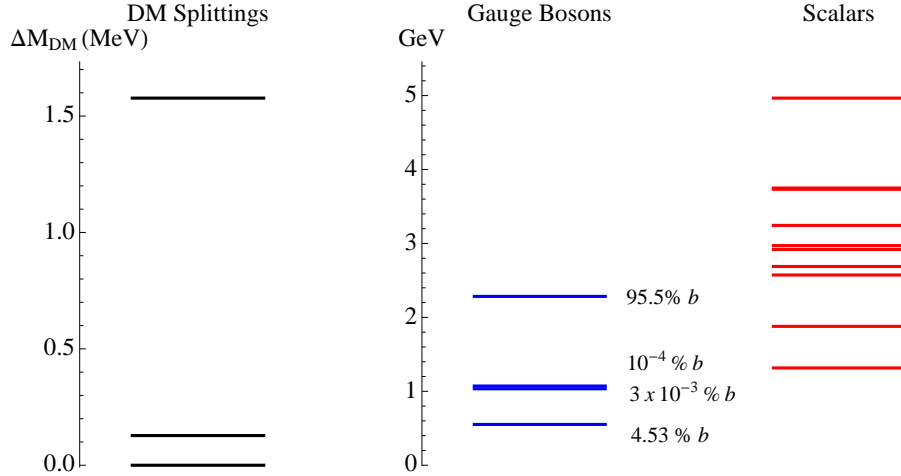

Figure 5: The spectrum of **Non-SUSY 1**, our two doublet non-SUSY benchmark. The left side shows the radiative mass splittings of the components of the dark matter triplet, measured from the ground state. The splittings allow for the XDM and iDM explanations of INTEGRAL and DAMA, respectively. The right side displays the spectrum of the GeV-scale dark sector. The b fractions of the gauge bosons are indicated and determine how strongly each gauge boson couples to Standard Model electromagnetic current. Because of custodial $SU(2)$, two of the gauge bosons are degenerate and do not mix with the b at tree-level, and these are the gauge bosons that couple between different dark matter states. They do mix with the b at one-loop, inducing a suppressed iDM coupling, and we include the dimension 6 operators $c_1^T |h_1 D h_1|^2$ and $c_2^T |h_2 D h_2|^2$ in order to parametrize custodial breaking corrections. The parameters of this benchmark are listed in the text.

5.1.2 Non-SUSY 2: Two Doublets and One Complex Triplet

We now add a complex triplet Higgs Φ to the two doublet model, with dark quantum numbers $\mathbf{3}_0$. The triplet vev breaks custodial symmetry, causing all gauge bosons to mix with the b and inducing the iDM coupling at tree-level. We take the triplet to be complex. While not the minimal possible choice, it has a more straightforward SUSY extension. We again choose a benchmark that breaks charge in the doublet sector and radiatively generates the XDM and iDM dark matter splittings. We have calculated its spectrum (Fig. 6).

For this benchmark, we choose the gauge couplings $g = 0.23$ and $g_y = 0.75$, and the dark matter hypercharge is chosen to be $y_{\rm dm} = 0.3$. The potential is similar to Eq. 2.6 except we take Φ to be complex:

$$
\begin{array}{c||c|c|c}
\text{field} & h_1 & h_2 & \Phi \\
\hline
\text{charge} & \mathbf{2}_{-1/2} & \mathbf{2}_{1/2} & \mathbf{3}_0
\end{array}
\tag{5.1}
$$

$$
V(h_1, h_2, \Phi) = V(h_1, h_2) + \frac{\lambda_\Phi}{2} \left(\text{Tr}\left[\Phi^\dagger \Phi \right] - |v_\Phi|^2 \right)^2 \tag{5.2}
$$
$$
+ \left(c_1 h_1^\dagger \Phi h_1 + c_2 h_2^\dagger \Phi h_2 + c_3 h_1^T \epsilon \Phi h_2 + \text{h.c.} \right)
$$

where the first term is the two doublet potential of Eq. 2.3. For the two doublet sector we choose the parameters: $v_1 = 1.8$ GeV, $v_2 = (1.8 + 1.4\,i)$ GeV, $\lambda_1 = 0.71$, $\lambda_2 = 0.47$, $\lambda_3 = 0.33$, $\lambda_4 = 0.099$, and $\cos\alpha = 0.052$. The parameters involving the triplet are chosen to be: $v_\Phi = (1.1 + 0.61\,i)$ GeV, $\lambda_\Phi = 0.51$, $c_1 = (0.054 + 0.47\,i)$ GeV, $c_2 = (0.74 + 0.69\,i)$ GeV, and $c_3 = (0.61 + 0.81\,i)$ GeV.

Figure 6: The spectrum of **Non-SUSY 2**, our two doublet and one complex triplet non-SUSY benchmark. The left side displays the radiative mass splittings of the dark matter triplet, measured from the ground state. The splittings can account for the XDM and iDM explanations of INTEGRAL and DAMA, respectively. The right side shows the dark sector spectrum, and the b fractions of the gauge bosons are indicated. The triplet vev breaks custodial $SU(2)$, and all 4 gauge boson mass eigenstates mix with the b at tree-level, although for this example most of the b is contained in two of the mass eigenstates. The parameters of this benchmark are listed in the text.

5.2 SUSY Benchmark Models

Now we consider two SUSY benchmarks, where the GeV scale is generated radiatively from interactions with the Standard Model, as discussed in Section 4. Since our models all employ a

kinetic mixing, they all receive SUSY breaking contributions from kinetic mixing mediation, as discussed in Section 4.1. In our first example, **SUSY 1**, this the only source of SUSY breaking, however as noted in appendix B, it is difficult within this framework to break charge with only one hypercharge neutral triplet. We circumvent this in this example by adding a second complex triplet and taking the triplets to have dark hypercharge. For our second example, **SUSY 2**, we add an additional gauge mediation source for GeV-scale SUSY breaking by taking dark matter to be charged as a $\mathbf{5} + \bar{\mathbf{5}}$ of the SM. Dark matter then acts as a messenger of gauge mediation, as discussed in section 4.2. For this setup, we can break charge with two doublets and one hypercharge neutral triplet.

5.2.1 SUSY 1: Kinetic Mixing Mediation with Two Doublets and Two Triplets

For this benchmark, we have two doublets, h_1 and h_2, and two complex triplets, Φ_1 and Φ_2, with dark quantum numbers $\mathbf{2}_{-1/2}$, $\mathbf{2}_{1/2}$, $\mathbf{3}_1$, and $\mathbf{3}_{-1}$. We have chosen triplet hypercharge assignments that allow Yukawa couplings between doublets and triplets, otherwise there may be pseudo-Goldstone bosons in the spectrum. The GeV scale is generated in the dark sector from kinetic mixing, as described in Section 4.1. The most general renormalizable superpotential for two doublets and two triplets with these charge assignments is:

$$
\begin{array}{c|c|c|c|c}
\text{field} & h_1 & h_2 & \Phi_1 & \Phi_2 \\
\hline
\text{charge} & \mathbf{2}_{-1/2} & \mathbf{2}_{1/2} & \mathbf{3}_1 & \mathbf{3}_{-1}
\end{array}
\tag{5.3}
$$

$$
W = \mu_h h_1^T \epsilon h_2 + \mu_\Phi \mathrm{Tr}\left[\Phi_1 \Phi_2\right] + \lambda_1 h_1^T \epsilon \Phi_1 h_1 + \lambda_2 h_2^T \epsilon \Phi_2 h_2
\tag{5.4}
$$

We include GeV scale μ and $B\mu$ terms for the doublets and triplets because they help break the dark gauge symmetry and lift runaway directions. We do not include the effects of running from the TeV scale to the GeV scale, which we take to be subdominant. Kinetic mixing mediation already gives negative scalar soft mass squareds at tree-level, leading to the breaking of dark gauge symmetry.

We include a kinetic mixing coefficient of $\epsilon = 2 \times 10^{-4}$, in terms of Eq. 2.1, which automatically generates the GeV scale in the dark sector. Our benchmark radiatively generates the XDM splitting, but unfortunately the smaller dark matter splitting is too large to account for iDM. We have calculated the mass spectrum (Fig. 7). The triplet vevs break custodial $SU(2)$ at tree-level and all gauge bosons mix with the b. The gauginos and Higgsinos are strongly mixed after dark symmetry breaking, but for this example the lightest fermion is a mostly gaugino-like Goldstino with a mass of only ~ 2 MeV. Such a field is present because SUSY is broken within the dark sector itself. The second lightest fermion, with mass ~ 190 MeV, is lighter than the lightest gauge boson. Thus, the dark gauge bosons will cascade into these light fermions, rather than SM lepton pairs. The second lightest fermion decays to the lightest fermion and a SM lepton pair through a 3-body decay, which can account for the astrophysical lepton production and lead to visibly displaced vertices at colliders. Another possibility, not realized in this example, is to have an approximately supersymmetric dark

sector, with a kinetic mixing mediation induced super-Higgs mechanism at a GeV. Gauginos then reside in massive vector supermultiplets and get GeV scale masses.

For this benchmark, we have chosen the gauge couplings $g = 0.22$ and $g_y = 1.2$, and dark matter hypercharge $y_{\rm dm} = 1/5$. The superpotential Yukawa couplings are chosen to be $\lambda_1 = 1.7 + 0.022\,i$ and $\lambda_2 = 0.5 + 1.8\,i$. For the doublets we choose $\mu_h = (0.11 + 0.63\,i)$ GeV and $(B\mu)_h = (0.74 + 0.69\,i)$ GeV^2. For the triplets we choose $\mu_\Phi = (0.51 + 0.83\,i)$ GeV and $(B\mu)_\Phi = (0.57 + 0.59\,i)$ GeV^2.

Figure 7: The spectrum of **SUSY 1**, our two doublet and two complex triplet SUSY benchmark with kinetic mixing mediation. The left side displays the radiative mass splittings of the dark matter triplet, measured from the ground state. The larger splitting allows for the XDM explanation of INTEGRAL, but the smaller splitting is too large to explain DAMA with iDM. The right side shows the dark sector spectrum, and the b fractions of the gauge bosons are indicated. The triplet vevs break custodial symmetry, and all four gauge boson mass eigenstates are part b at tree-level. The gauginos and Higgsinos are strongly mixed after dark symmetry breaking. The lightest fermion, with mass ~ 2 MeV, is mostly gaugino and light because gauginos get no soft masses from kinetic mixing mediation. The second lightest fermion has a 3-body decay to the lightest fermion and a SM lepton pair, which can account for astrophysical lepton production and lead to a visibly displaced vertex at colliders.

5.2.2 SUSY 2: Gauge Mediation with Two Doublets and One Triplet

For this benchmark, we supersymmetrize the Higgs content of our **Non-SUSY 2** benchmark, including SUSY breaking contributions from both kinetic mixing mediation and gauge mediation with dark matter messengers. The most general renormalizable superpotential for two doublets and one triplet is the following:

$$\begin{array}{c|c|c|c}
\text{field} & h_1 & h_2 & \Phi \\\hline
\text{charge} & \mathbf{2}_{-1/2} & \mathbf{2}_{1/2} & \mathbf{3}_0
\end{array} \tag{5.5}$$

$$W = \mu_h h_1^T \epsilon h_2 + \mu_\Phi \text{Tr}\left[\Phi^2\right] + \lambda h_1^T \epsilon \Phi h_2 \tag{5.6}$$

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As in the **SUSY 1** benchmark, we include GeV scale μ terms for the doublets and triplet and do not include the effects of running from the TeV scale to the GeV scale. There are already negative scalar soft mass squareds at tree-level because of the nonzero dark matter supertrace [38], leading to the breaking of dark gauge symmetry. For this example, it is not necessary to include GeV scale $B\mu$ terms for the doublets or triplet.

We have chosen a benchmark which generates a GeV scale dark sector with charge breaking and custodial breaking, and which leads to radiative XDM and iDM splittings. We have calculated the mass spectrum (Fig. 8). The gauginos and Higgsinos are strongly mixed after dark symmetry breaking, but the three heaviest fermions with masses near ~ 5.5 GeV are almost pure Higgsino mixtures. The spectrum is slightly split by a small separation between the dark μ and soft mass scales. The gauge couplings are chosen to be $g = 0.3$ and $g_y = 0.37$ and the dark matter hypercharge is $y_{\rm dm} = 1/2$. The kinetic mixing is $\epsilon = 7 \times 10^{-5}$ in terms of equation 2.1. The messenger scale of SUSY breaking to the Standard Model, which enters the log divergence of Eq. 4.8, is chosen to be $\Lambda_{UV} = 30$ TeV, corresponding to low-scale gauge mediation. The standard model doublet dark matter mass components, in terms of Eq. 4.6, are given by $m_f^{(\mathbf{2})} = 800$ GeV, $m_s^{(\mathbf{2})} = 50$ GeV, and $B\mu^{(\mathbf{2})} = (40 \text{ GeV})^2$. As discussed in section 4.2, the small soft mass is needed to generate the GeV scale in the dark sector, and we choose a small $B\mu$, keeping dark matter fermionic. For the standard model triplet components of the dark matter $\mathbf{5} + \bar{\mathbf{5}}$, we choose the parameters $m_f^{(\mathbf{3})} = 840$ GeV, $m_s^{(\mathbf{3})} = 50$ GeV, and $B\mu^{(\mathbf{3})} = (300 \text{ GeV})^2$. The larger $B\mu$ for the triplet leads to GeV scale gaugino soft masses in the dark sector (see Eq. 4.9 and the surrounding discussion). The superpotential parameters are $\mu_h = (0.27 + 0.28\,i)$ GeV, $\mu_\Phi = (2.52 + 3.48\,i)$ GeV, and $\lambda = 0.29 + 1.51\,i$.

6. Signals of a Non-Abelian Dark Sector at the Tevatron and the LHC

In this section we discuss the collider phenomenology associated with the models presented in the previous sections. In the first part of this section we analyze the generic predictions associated with a non-abelian dark sector that is linked to the SM only via kinetic mixing. In the second part we present the signals expected in supersymmetric versions of such models. Throughout, we limit the discussion to the Tevatron and LHC. It is important to realize that in the case of a GeV scale dark sector such high-energy accelerators are needed more for their luminosity than their energy reach, which is considerably higher than the dark sector scale. In supersymmetric implementations, colored MSSM superpartners can be copiously produced at hadron colliders. Their subsequent decay into dark states can produce spectacular signals involving multiple lepton jets. We leave it for future work to investigate the phenomenology of these models at low-energy experiments, but see Ref. [11] for low-energy signatures of similar models.

A new sector of light particles with very weak couplings to the Standard Model have been discussed in detail in the context of the "Hidden Valley" models [40]. Their collider phenomenology was investigated in [41, 42]. In particular, the modifications such models can introduce to the decay chains of the MSSM was clarified in Ref. [43]. Here, we focus on the
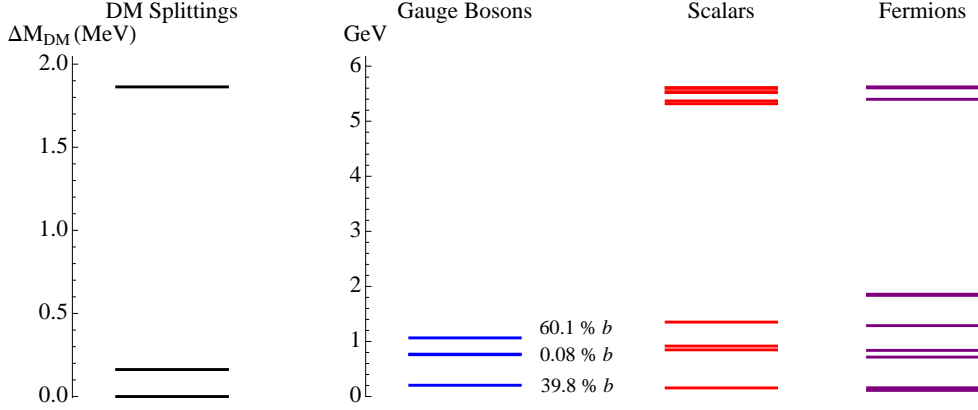
Figure 8: The spectrum of **SUSY 2**, our two doublet and one complex triplet SUSY benchmark with both kinetic mixing mediation and dark matter messengers charged under the Standard Model. The left side displays the radiative mass splittings of the dark matter triplet, measured from the ground state. The splittings can allow for the XDM and iDM explanations of INTEGRAL and DAMA, respectively. The right side shows the dark sector spectrum, and the b fractions of the gauge bosons are indicated. The triplet vev breaks custodial $SU(2)$ and three of the gauge boson mass eigenstates are part b at tree-level. The dark spectrum now includes GeV scale fermions, and the gauginos and Higgsinos are strongly mixed after dark symmetry breaking. The three heaviest fermions, with masses near ~ 5.5 GeV, are almost pure Higgsino mixtures. The spectrum is slightly split by a small separation between the dark μ and soft mass scales. The parameters of this benchmark are listed in the text.

particular scenario which uses the kinetic mixing as the essential link between the SM and the dark sector. In addition, motivated by astrophysical observations, we allow the dark sector to decay back to light leptons ($e^\pm$ and $\mu^\pm$) only. For the purpose of this paper, we do not concern ourselves with a possibly small branching fraction into pions.

The pair production of the dark matter states at colliders is certainly possible if they happen to carry SM weak charge. However, their detection proves extremely difficult since they are not accompanied by any hard object. Even if the excited states of its SM multiplet are produced, their decays are too soft to trigger on since they are separated by only $\sim$ GeV (notice that this splitting is generated by the SM gauge interactions and are of order αM_Z [34]).

6.1 Production and decay of dark gauge bosons and Higgses

As discussed in detail in appendix A, the kinetic mixing induces two important, ϵ suppressed, couplings: The SM electromagnetic current is now also charged under the dark gauge bosons; the SM Z^0 boson is now coupled to the dark hypercharge current. Before discussing each of these couplings and their impact on collider signals, let us briefly discuss the decays of the dark gauge boson and the dark Higgses.

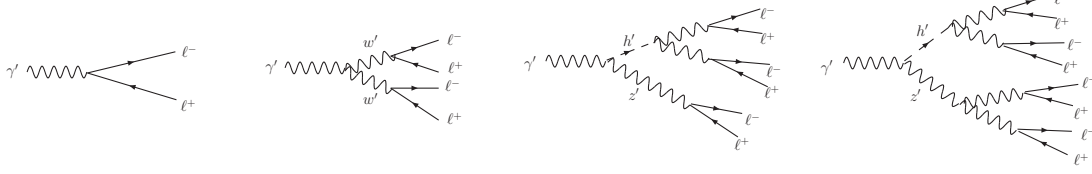### 6.1.1 Dark gauge boson and Higgs decay chains

Figure 9: Typical decay chains starting with a dark gauge boson, labelled γ' in this plot. The dark decay chain can have several stages and involve additional dark sector states, such as other dark gauge bosons (labeled w', z' in this figure), and dark Higgses (labelled h').

The non-abelian nature of the dark sector implies the presence of complicated decay chains. Some of the typical decays chains are shown in Fig. 9. In the dark sector, gauge boson mass eigenstates are generically mixtures of all four $SU(2)_{\mathrm{dark}} \times U(1)_y$ gauge eigenstates. In Fig. 9 and the rest of this section, we have used γ' (and also w' and z' in this figure) to denote any one of these mass eigenstates. For an abelian dark sector with kinetic mixing with the SM, γ' decay leads to a di-lepton final state, shown in the first panel from the left of Fig. 9. On the other hand, a non-abelian dark sector, like one of the examples considered in this paper, leads to complicated decay chains, such as the ones shown in the rest of Fig. 9. The dark Higgs sector, necessary to break the non-abelian group, may also participate in such cascades as shown in the right two panels of Fig. 9. Such cascades inevitably produce multiple, easily > 2 and possibly 8, final state leptons, which provides a unique signature of the non-abelian nature of the dark sector[11]. We expect the decay between dark states to be generically prompt. Therefore, the decay length is dominated by the very last decays back into SM leptons. A rough estimate for a generic decay is then,

$$c\tau_{2-\mathrm{body}}^{\gamma'\rightarrow n\ell} \sim \frac{1}{\alpha\epsilon^2 m_{\gamma'}} = 2.7 \times 10^{-6} \ \mathrm{cm} \ \left(\frac{\mathrm{GeV}}{m_{\gamma'}}\right)\left(\frac{10^{-3}}{\epsilon}\right)^2. \tag{6.1}$$

With moderate boost $\gamma \sim \mathcal{O}(10)$, this may lead to a displaced vertex if $\epsilon \lesssim 10^{-4}$.

To be observable at hadron colliders, the dark boson which initiates such a cascade must carry $p_{\mathrm{T}} \sim \mathcal{O}(10s)$ GeV. Therefore, regardless of the precise nature of the cascade which ensues, its decay products have small opening angles $\delta\theta \sim m_{\gamma'}/p_{\mathrm{T}} < 0.1$. Those decay products will eventually decay into several collimated SM leptons. A collection of more than 2 hard and collimated leptons is dubbed a "lepton jet" [28].

6.1.2 Displaced vertices and missing energy

While Eq. (6.1) is the generic estimate for the resulting decay length of dark cascades, there are several exceptions which may result in more noticeably displaced vertices or missing energy in lepton jets.

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[11]Sometimes phase space constrains the flavor of the lepton. For example, a GeV dark gauge boson cannot decay into more than 4 muons

If it is kinematically forbidden for a dark gauge boson to have 2-body on-shell decays within the dark sector, then the dark gauge boson may decay directly into two leptons. However, a noticeable exception occurs when the 3-body decay $\gamma' \to a'^* b_1' \to b_1' b_2' b_3'$ is kinematically allowed, where a', b' can be either dark gauge boson or dark Higgs states. In this case, there is an additional suppression of $(\delta m/m_{a'})^5 \times (\text{3-body phase space})$ on the decay width, where $\delta m \sim m_{\gamma'} - \sum_i m_{b_i'}$, and we have used $m_{\gamma'} \sim m_{a'}$ in this estimate. This decay channel can be competitive and even dominate over the direct decay into 2 leptons. In particular, when the decay into SM leptons is strongly suppressed (dark pseudoscalar decay) or all together forbidden (dark fermion decay), the 3-body process may dominate and lead to a displaced vertex. The impact parameters of multiple leptons associated with this displaced vertex will not be correlated with each other since they come from the decays of different resonances $b_{1,2,3}'$.

If the lightest dark sector state is a dark Higgs, h_0', it cannot directly decay into SM leptons (unless it mixes the SM Higgs, see Ref. [19, 1]). In this case, the dark Higgs will either decay into 4 leptons through two off-shell gauge bosons, shown in the left panel of Fig. 10, or into 2 leptons through a one-loop decay. Either way, such a decay leads to a very long life-time, $c\tau \sim \mathcal{O}(\text{km})$ for $m_{h_0'} \lesssim$ GeV. In this case, dark cascades which involve this lightest scalar contain missing energy as it escapes the detector. These cascades can still produce observable lepton jets because, in addition to missing energy, one still gets leptons from the intermediate steps of the decay, such as $h_i' \to a' h_0'$ followed by $a' \to$ lepton pairs. In this case, the lepton jet contains missing energy that is collimated with the leptons of the same cascade.
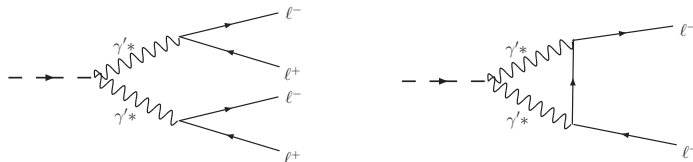

Figure 10: Two possible decay channels if the lightest dark sector state is a scalar from the dark Higgs sector.

An additional source of missing energy comes in a supersymmetric dark sector with R-parity. The lightest dark supersymmetric particle (LDSP) may be stable if the gravitino is heavier. Otherwise, it may eventually decay into the gravitino. Either way, it carries with it missing energy. Unless the MSSM sector decays directly into the LDSP, in which case there may be no lepton jets, missing energy due to the LDSP will be collimated with the visible lepton jets, very similarly to the non-SUSY case. Such correlations provide an additional handle on the reconstruction of these events since we know the direction of the missing particles and can treat them as having vanishing masses. We provide an example of such a reconstruction in the case of rare Z^0 decays below.

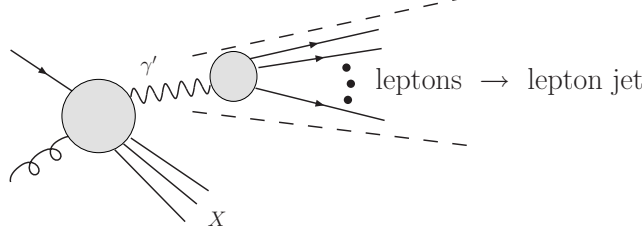### 6.1.3 Direct Production

Figure 11: Direct production of a dark gauge boson in a process very similar to prompt photon production in the Standard Model.
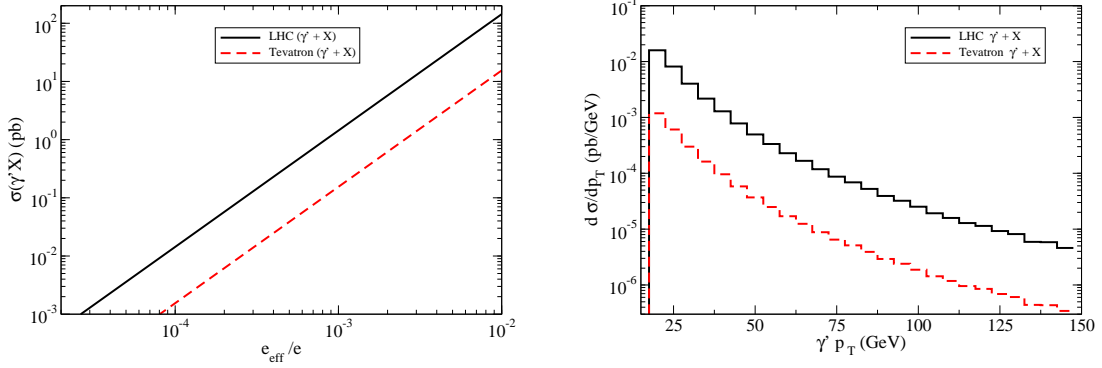

Figure 12: In the left pane, we show the rate of direct production of the dark gauge boson as a function of e_{eff}/e, where e_{eff} is the effective coupling of dark gauge boson to the Standard Model fields.

The kinetic mixing between the dark force carrier and the SM photon induces a small dark charge for electromagnetically charged SM fields. Consequently, the dark gauge boson can be directly produced in colliders via a process analogous to prompt photon production in the SM, shown in Fig. 11.

In the left panel of Fig. 12, we present the production rate of dark gauge bosons as a function of e_{eff}/e, where $e_{\text{eff}} = \epsilon e \cos \theta_W f_b$ is their effective gauge coupling to SM fields[12] and f_b is the fraction of the dark hypercharge gauge boson b_μ in a given dark gauge boson mass eigenstate. In the right panel of Fig. 12 we plot the inclusive differential cross-section of dark photon (γ') production at the LHC and the Tevatron with $e_{\text{eff}} = 10^{-3}e$.

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[12]The simulation was actually of prompt photon production with PYTHIA [44] and the resulting cross-section was multiplied by a factor of e_{eff}^2/e^2.
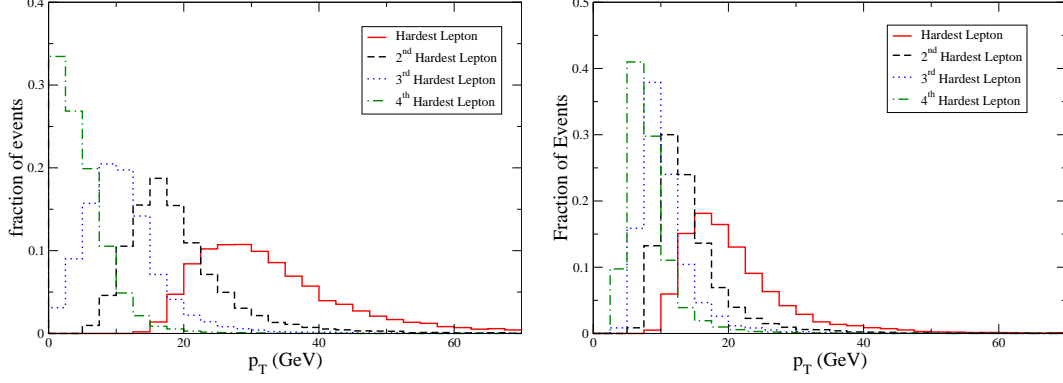
Figure 13: p_T distributions for cascades resulting in 4 (left) and 8 (right) leptons, for events with $p_T > 50$ GeV for γ'.
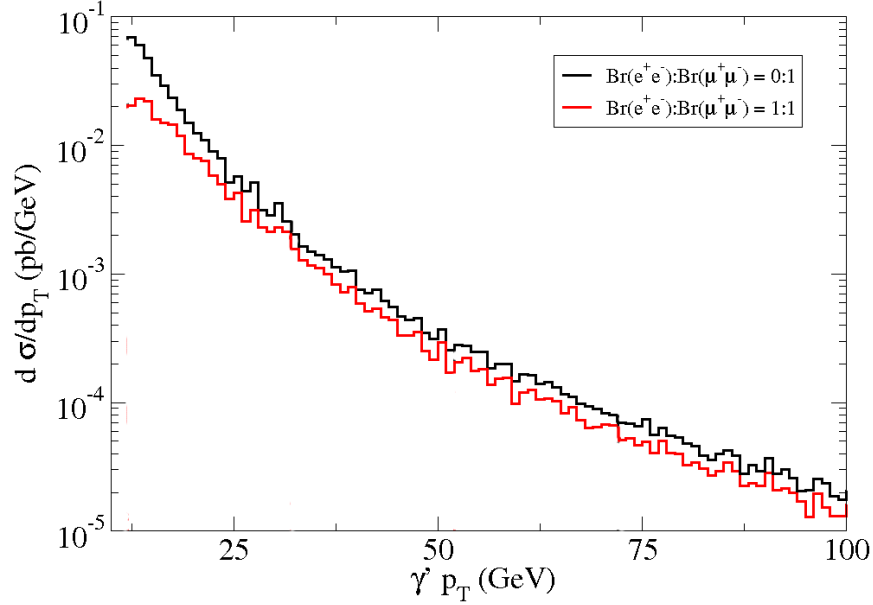

Figure 14: The differential cross-section as a function of the p_T of γ' at the LHC ($\sqrt{s} = 14$ TeV) after including muon triggers, demanding either a single muon with $p_T > 7$ GeV or two muons with $p_T > 3$ GeV. Proper η cuts were imposed and each event was required to contain at least 3 leptons.

After dark vector bosons are produced, they typically cascade down to multiple leptons that form a lepton jet as discussed above. These leptons carry a significant amount of p_T, as shown in Fig. 13. At CMS, the Level 1 Dimuon trigger (2 muons with $p_T > 3$ GeV in $|\eta| < 2.4$) or single muon trigger (1 muon with $p_T > 7$ GeV and $|\eta| < 2.4$) should be able to detect those events that contain muons [45]. The electron triggers are single e (isolated $E_T > 26$ GeV), double e (isolated $E_T > 14.5$) and double relaxed e (not isolated $E_T > 21.8$ GeV). Since the
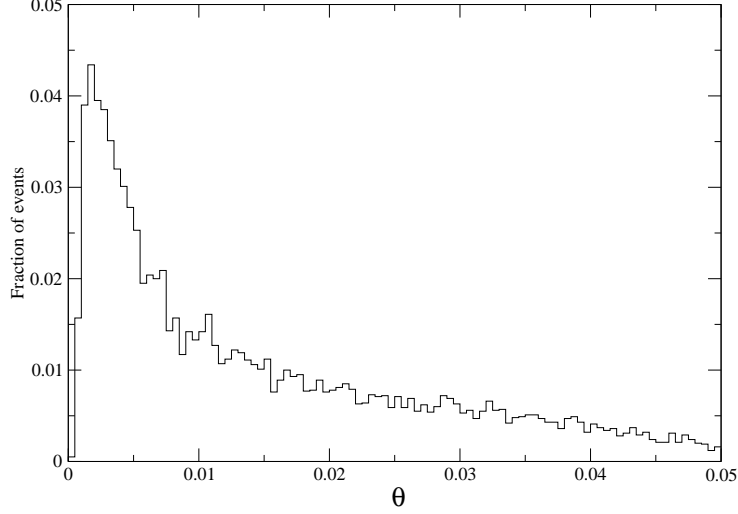
25

Figure 15: The fraction of events with a maximum opening angle θ between leptons in a lepton jet, which contains 4 leptons.

resulting electrons are unlikely to be isolated "electromagnetic" objects, the double relaxed e is probably necessary. We will conservatively assume that muons alone are triggered on. In Fig. 14, we show the differential cross section of dark γ', taking into account the simple requirements on muon triggering.

6.1.4 Distinguishing Leptons

Let us discuss the issue of discriminating individual leptons within a given lepton jet[13]. In our present discussion, we focus on muons. In Fig. 15 we plot the maximal opening angle between any two of the four leptons. At such high momenta, the resulting decay products are highly collimated with an initial opening angle of approximately $\theta \sim m_{\gamma'}/p_T < 0.1$, which can be as small as 10^{-2}. By the time these muons reach the first layer of the muon system, they typically acquire a sufficient separation to be distinguished. For example, as depicted in Fig. 13, a typical scenario will have two muons with average $p_T \sim 20$ GeV and $\Delta p_T \sim 5$ GeV. Without even including the initial lepton jet opening angle, we estimate that the acquired separation is about 10 cm (in the CMS detector), which is greater than the cell size of ~ 4.5 cm. The separation between two same sign muons is proportional to $\Delta p_T/(p_T^{\mu_1} p_T^{\mu_2})$. For a given lepton jet p_T, since both Δp_T and p_T are inversely proportional to the number of leptons, higher multiplicities actually result in larger separations. We also notice from Fig. 13 that leptons in lepton jets typically have different p_T, such that $\Delta p_T/p_T \sim 20\%$ or so. LHC detectors can achieve better muon momentum resolution. For example, CMS can achieve $\Delta p_T/p_T \sim 1\%$ (about 10% with muon system only) in the momentum regime of interest

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[13]We are grateful for valuable discussions with Jim Olsen on this subject.

[46]. ATLAS can achieve a similar precision [47]. Finally, the muon isolation separation defined by CMS can be as small as $\Delta R = 0.01$. The angular resolution is even smaller, about 2 mrad[46][14]. Thus, it is reasonable to assume that CMS will be capable of resolving several, if not all of the muons. The primary background arises from K and π decays and $J/\psi \to \mu^+\mu^-$ (muons coming off soft jets can be vetoed with isolation cuts), and possibly from other heavy flavor decays. The high lepton multiplicity in those events and the lack of hadronic activity around the lepton jet should be sufficient to fight these backgrounds and obtain a clean sample. However, a more careful collider analysis is certainly warranted, but is beyond the scope of the present work.

6.1.5 Rare Z^0 decay

As discussed in Appendix A, the kinetic mixing also induces a coupling $\epsilon Z_\mu J_b^\mu$, where J_b^μ is the dark hypercharge current. Thus, we can produce dark hypercharged states through rare decays of the Z^0, shown schematically in Fig. 16. The ϵ^2 suppression makes LEP searches irrelevant due to luminosity limits, but the Tevatron and LHC may probe such events. The decay branching ratio to any particular dark sector state d_i can be written as

$$\text{BR}(Z^0 \to d_i d_i) = \frac{c_{d_i}}{\Gamma_Z^0} \frac{\epsilon^2 g_y^2 y_{d_i}^2 \sin^2 \theta_W}{48\pi} M_{Z^0}, \tag{6.2}$$

where c_{d_i} depends on decay matrix element and is proportional to the number of degrees of freedom of d_i. The total branching ratio into the dark sector will scale linearly with the number of dark sector states, which could be easily $\mathcal{O}(10)$ in our scenario.
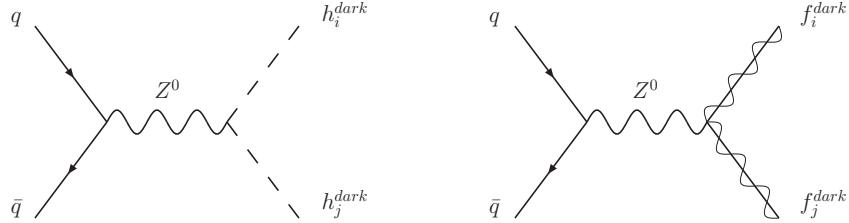

Figure 16: Z^0 production and two possible decay channels into the dark sector. On the left we depict a decay into the dark Higgs sector. Fermionic channels, such as the one shown on the right, are dominantly associated with the Higgsino states possible in supersymmetric versions of the model.

The SM photon does not couple to the dark sector states. However, there is a "continuum" contribution to the same amplitude through off-shell dark photon, $q\bar{q} \to \gamma'^\star \to d_i d_i$, which is proportional to $e_{\text{eff}}^2 \propto \epsilon^2$. Depending on the spectrum and couplings in the dark sector, it could have important contributions to the signal off the Z-peak. In this section, we will focus on the contribution within the Z resonance.

---

[14]This is the resolution quoted for a single hit. It is beyond our abilities to evaluate how the resolution deteriorates with multiple collimated muons.

The production rates of dark sector states at the Tevatron and LHC are shown in Fig. 17 [48]. We present rates coming from decays into bosonic (denoted by h') and fermionic (denoted by f') dark sector states. In the context of the SUSY models discussed later in this section, these bosonic and fermionic states could refer to dark Higgs bosons or Higgsinos, respectively. On the other hand, the collider phenomenology is similar if other possible dark sector states decay into lepton jets. A cut of $|\eta| < 2.4$ has been imposed on the direction of the lepton jets. The difference in rates between the fermionic and bosonic channels results from the η cut, the boost of the Z^0 in the lab frame, and the fact that fermions are more likely to be emitted along the boost direction because of angular momentum conservation.

Figure 17: Left: The production rate as a function of the branching ratio of the decay: $Z^0 \rightarrow$ dark states. The solid and dashed lines are for Z^0 decays into dark sector scalars and fermions, respectively. Right: lepton jet p_T distribution resulting from Z^0 decays.

As can be seen from Fig. 17, the lepton jets produced in this way are peaked towards $p_{\mathrm{T}}^{\text{lepton jet}} \sim 0.5\, M_Z$. Therefore, they are typically harder than the lepton jets resulting from the prompt production of dark gauge bosons. As we have discussed in Section 6.1.3, such harder lepton jets will be easier to trigger on. However, we expect the efficiency of identifying different leptons in a lepton jet will be lowered as it is $\propto 1/p_T$.

Reconstructing the Z^0 is not difficult and helps to reduce the background. With enough statistics, it is even possible to study the angular distribution of the resulting lepton jets and get a handle on the spin of the dark sector states as demonstrated in Fig. 18. About 5000 events are used in this plot. We see that the expected rise for $\cos\theta \sim \pm 1$ from the fermionic decay channels is washed out due to the large boost of Z^0 and the $|\eta|$ cut. However, the resulting distribution is still quite different from that of the bosonic decay channel.

6.2 Collider signals of supersymmetric models

In this section, we discuss the collider signals associated with supersymmetric models. In
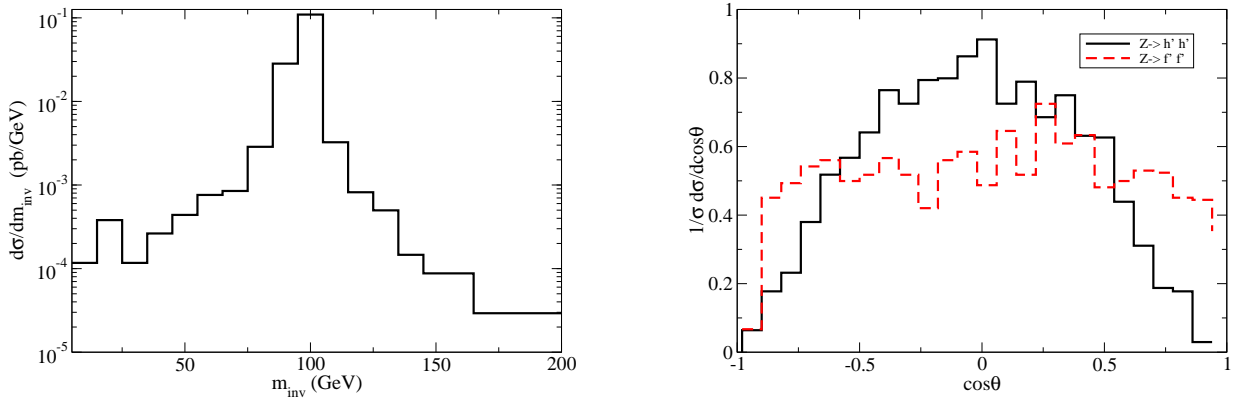
Figure 18: Left: Reconstructed Z^0 boson. Right: Normalized lepton jet angular distribution in the Z^0 boson's rest-frame.

Section 4, we have focused on models with low supersymmetry breaking scale. However, the present discussion of the resulting collider signatures is largely independent of that scale or other MSSM details since we will not consider any specific superpartner spectrum. In that sense, models with higher supersymmetry breaking scales, such as the Planck slop option suggested in Ref. [28], are only different in that the gravitino is heavier. Hence, the end of the dark sector decay chain will not involve the gravitino. However, this does not have a visible effect on the collider signals. Even in the low scale models where the gravitino is light, the decay length of the dark sector LSP to the gravitino is much larger than the detector size. That said, it is important to note that the collider signatures discussed in this section are based on the assumption that the MSSM LSP dominantly decays into the dark sector.

With supersymmetry, the dark sector states are dominantly produced from cascade decays of MSSM colored superpartners, such as gluinos and squarks. These particles follow typical MSSM decay chains down to the MSSM LSP (not the gravitino). The effect of the GeV dark sector is to extend and/or modify the decay chains following MSSM LSP production [28]. We begin by summarizing the main features of such cascades.

Let us first note, however, that a notable exception occurs if dark matter is part of a pair of $\mathbf{5} + \bar{\mathbf{5}}$ under the SM gauge groups. An example of such a model was presented in the benchmark of Section 5.2.2. The rate for the production of the colored components of such a pair is shown in Fig. 19. Thus, the LHC has great potential for producing such states up to about 2 TeV. As pointed out in Ref. [28], as long as the colored particles decay only through higher dimensional operators they will be long-lived and may have decays with very distinct signatures [49]. We will not elaborate on these possibilities but refer the interested reader to the detailed discussion in Ref. [28].
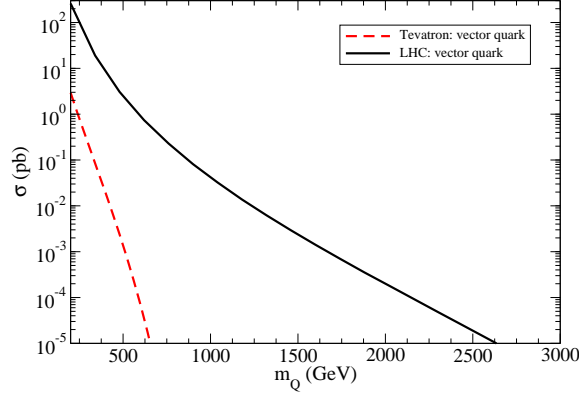
Figure 19: Production rate of one set of $3 + \bar{3}$ vector-like quarks, which can be part of the dark matter multiplet. The rate should be scaled by the number of such representations and the dimension of the dark matter representation under G_{dark}.

6.2.1 MSSM decays into the dark sector

Kinetic mixing implies that if the MSSM LSP is a neutralino then it decays into dark sector states with a lifetime of

$$
\begin{aligned}
\tau_{\text{LSP}\to h+\tilde{h}} &\sim \left(\alpha_y^{\text{dark}} f_{\tilde{B}}^2 \epsilon^2 M_{\text{LSP}} \right)^{-1} \\
&= 7 \times 10^{-19} \text{ s } \left(\frac{100 \text{ GeV}}{M_{\text{LSP}}} \right)^2 \left(\frac{0.01}{\alpha_y^{\text{dark}}} \right) \left(\frac{1.0}{f_{\tilde{B}}} \right)^2 \left(\frac{10^{-3}}{\epsilon} \right)^2,
\end{aligned}
\tag{6.3}
$$

where $f_{\tilde{B}}$ is the bino fraction of the MSSM LSP. In the low-scale gauge mediation models constructed earlier in this paper, it is possible for the gravitino to be significantly lighter than the MSSM LSP. When the gravitino is light, another possible decay channel for the MSSM LSP is LSP$\to X_{\text{SM}}\tilde{G}$, where X_{SM} can be a photon, Z, or Higgs, depending on the model parameters and phase space. The decay lifetime can be estimated as

$$
\tau_{\text{LSP}\to\gamma,Z,h+\tilde{G}} \sim \left(\frac{M_{\text{LSP}}^5}{16\pi F^2} \right)^{-1} = 3.3 \times 10^{-13} \text{ s } \left(\frac{100 \text{ GeV}}{M_{\text{LSP}}} \right)^5 \left(\frac{\sqrt{F}}{100 \text{ TeV}} \right)^4.
\tag{6.4}
$$

We see that the LSP dominantly decays into the dark sector instead of the gravitino. However, the two channels can be comparable in certain regions of parameter space, such as $f_{\tilde{B}} \sim 0.1$ and a low supersymmetry breaking scale $\sqrt{F} \sim \sqrt{m_{3/2}M_P} \sim 10$ TeV.

When the MSSM LSP is a sfermion ($\tilde{\ell}$ or $\tilde{q}$), things become more subtle. One possible decay channel is through an off-shell gaugino with a significant bino fraction, $\tilde{f} \to f + \tilde{\chi}^* \to f + [\text{dark sector states}]$, shown the left panel of Fig. 20. Its decay lifetime can be estimated
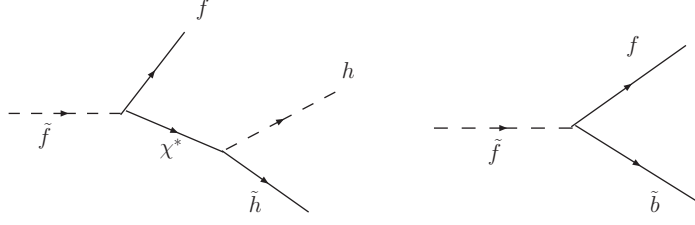
Figure 20: Decay of sfermion LSP.

to be

$$\tau_{\tilde{f}\to 3-\text{body}} \sim \left[\alpha_y^{\text{dark}} g_Y^2 c_{f\chi}^2 f_{\tilde{B}}^2 \epsilon^2 \frac{m_{\tilde{f}}}{16\pi^2} P(m_{\tilde{f}}/M_\chi) \right]^{-1} \tag{6.5}$$

$$= 8.3 \times 10^{-16} \text{ s} \left(\frac{100 \text{ GeV}}{m_{\tilde{f}}} \right) \left(\frac{0.01}{\alpha_y^{\text{dark}}} \right) \left(\frac{1.0}{c_{f\chi} f_{\tilde{B}}} \right)^2 \left(\frac{10^{-3}}{\epsilon} \right)^2 \frac{1}{P(m_{\tilde{f}}/M_\chi)},$$

where $c_{f\chi}$ is the effective fermion-χ coupling, and $P(m_{\tilde{f}}/M_\chi)$ is a function with the limit $P \to (m_{\tilde{f}}/M_\chi)^4$ for $M_\chi \gg m_{\tilde{f}}$. Another possible decay channel is $\tilde{f} \to f + \tilde{b}$, shown in the right panel of Fig. 20, where $\tilde{b}$ is the dark bino. However, as explained in Appendix A, in addition to kinetic mixing, this coupling has an additional suppression of order $M_{\tilde{b}}/M_{\tilde{B}}$. Hence, its lifetime is

$$\tau_{\tilde{f}\to f+\tilde{b}} \sim \left[\alpha_Y \epsilon^2 m_{\tilde{f}} \left(\frac{M_{\tilde{b}}}{M_{\tilde{B}}} \right)^2 \right]^{-1}$$

$$= 6.6 \times 10^{-15} \text{ s} \left(\frac{100 \text{ GeV}}{m_{\tilde{f}}} \right) \left(\frac{10^{-3}}{\epsilon} \right)^2 \left(\frac{1 \text{ GeV}}{M_{\tilde{b}}} \right)^2 \left(\frac{M_{\tilde{B}}}{100 \text{ GeV}} \right)^2. \tag{6.6}$$

Notice that when the off-shell gaugino state is dominantly bino, we have

$$\frac{\tau_{\tilde{f}\to 3-\text{body}}}{\tau_{\tilde{f}\to f+\tilde{b}}} \sim \frac{4\pi}{\alpha_y^{\text{dark}}} \frac{M_{\tilde{b}}^2 M_{\tilde{B}}^2}{m_{\tilde{f}}^4}. \tag{6.7}$$

Therefore, these two channels can be either quite different or comparable, depending very sensitively on the details of the model. In principle, these two decay channels are distinguishable experimentally, as the three (two) body decay gives rise to three (two) different lepton jets, respectively. Notice also that in this case, it is easier for the channel that decays into the gravitino to be competitive as well, if F is close to tens of TeV.

Once the cascade has progressed into the dark sector, it will decay through the mass hierarchy there. Several typical cascades in the supersymmetric dark sector are shown in Fig. 21. For example, the center panel on the second row is important when the fermionic dark superpartners are lighter than the dark gauge boson. An example of such a scenario is
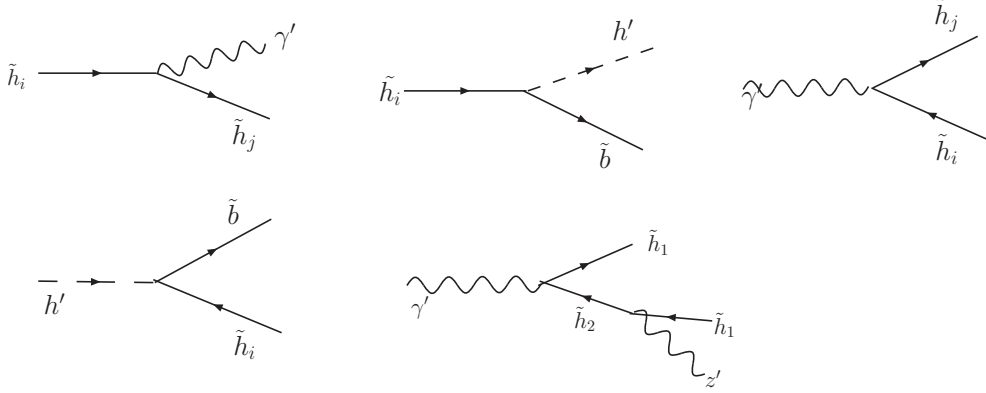
Figure 21: Typical SUSY dark sector decay chains. The dark sector states, γ', h' and z', can be either on or off shell. They will cascade further to produce lepton jets, similarly to the non-SUSY case.

discussed in Section 5.2.1. All of the decay products in the same chain are collimated into one lepton jet with typical $p_T \sim 100$ GeV. Therefore, it is difficult to uncover the details of the decay chain that produces a given lepton jet.

Finally, we discuss the endpoint of the dark sector decay chain. First, consider the situation where the lightest dark sector particle is the LDSP, as in the benchmarks of sections 5.2.1 and 5.2.1. If the gravitino is heavier than the LDSP, then the decay chain will end there with the LDSP escaping the detector, producing missing energy along the direction of the lepton jet. If the gravitino is lighter than the LDSP, as in the case of low scale supersymmetry breaking models, the last step of the cascade will be LDSP$\to X_{\mathrm{SM}}\tilde{G}$, where X_{SM} corresponds to a light SM particle, such as a photon or lepton. For simplicity, we consider the case where the LDSB is mostly dark bino. We can estimate its decay lifetime as

$$\tau_{\tilde{b}\to\gamma\tilde{G}} \sim \left[\frac{\epsilon^2}{16\pi}\frac{M_{\tilde{b}}^5}{F^2}\right]^{-1} = 3.3 \times 10^3 \text{ s } \left(\frac{10^{-3}}{\epsilon}\right)^2 \left(\frac{1 \text{ GeV}}{M_{\tilde{b}}}\right)^5 \left(\frac{\sqrt{F}}{100 \text{ TeV}}\right)^4, \qquad (6.8)$$

which is clearly not relevant on the collider timescale. Second, we consider the case where there is also a dark sector gauge boson, b, that is lighter than the LDSP. This situation is not realized in our benchmarks of Section 5.2, but is certainly a possibility. In this case, the LDSP can decay through the channel, $\tilde{b} \to b\tilde{G}$, where b subsequently decays to leptons. There is no ϵ^2 suppression here, but setting ϵ to 1 in Eq. (6.8) gives a decay length inside the detector only for $\sqrt{F} \lesssim 10$ TeV. Therefore, we can effectively think of the LDSP as the endpoint of the dark sector decay chain for most of parameter space.

6.2.2 Extended discovery reach for direct electroweak-ino production

The direct production of electroweak-inos is an important channel since it is independent of the existence of colored superpartners and may provide additional information on the
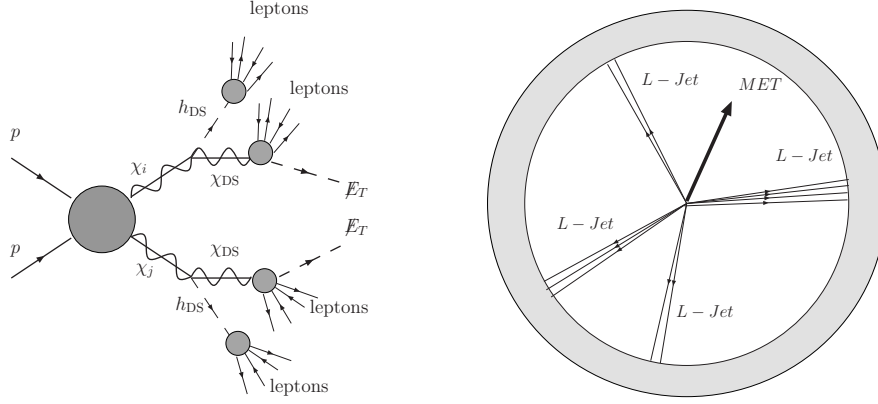
Figure 22: Pair-production of the SM LSP can result in spectacular lepton jet + MET events. On the left we depict the event topology. On the right we show a schematic representation of the resulting geometry.
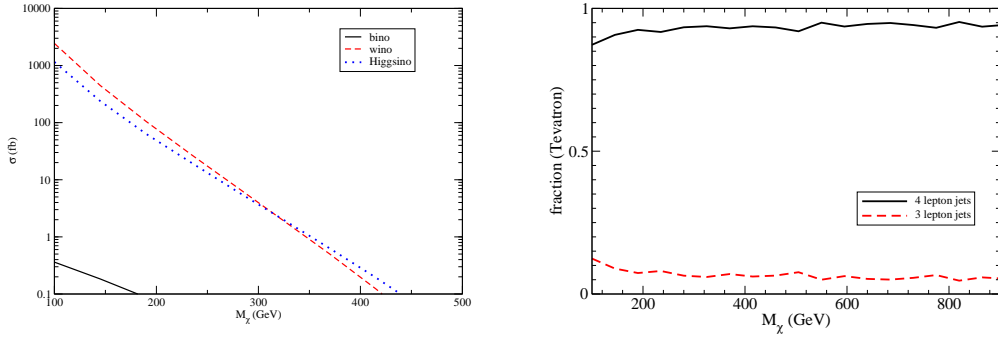

Figure 23: Left: The cross-sections for electroweak-ino production at the Tevatron. We have included both LSP pair production and, in the case of wino and Higgsino LSP, the production of closely degenerate states, as a function of M_χ. We choose the squark mass to be 750 GeV. Right: the fraction of events with 3 and 4 lepton jets within the central region $|\eta| < 2.4$.

properties of those electroweak-inos. In the conventional MSSM, it is usually difficult to see events with direct pair-production of electroweak-inos. In the case of direct MSSM LSP production, one has to trigger on some additional hard radiation, which has a lower rate and a large background. The pair-production of heavier electroweak-ino states which cascade down to the LSP may be easier to observe but suffers from large SM background. However, in the scenario we consider, the LSP of the MSSM, which we denote by χ_0, will decay further into the dark states [43] whose decays result in leptons and missing energy [28]. Such events are easy to trigger on since all the leptons carry significant amounts of p_T. Since χ_0 is produced almost on threshold, its boost factor is order unity and the opening angle in the
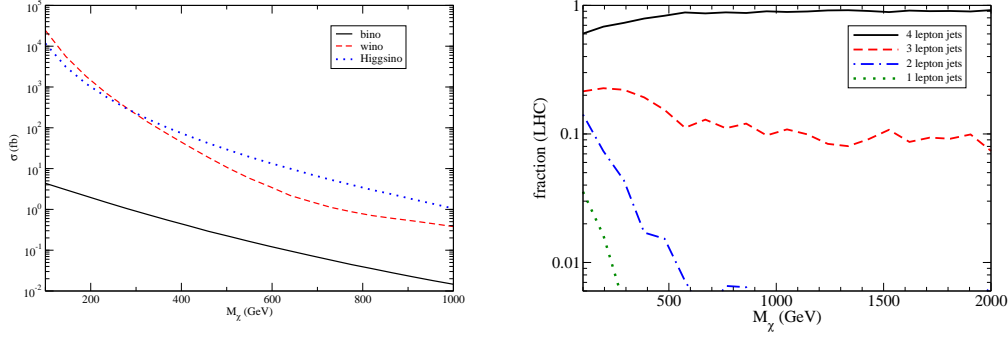
Figure 24: Left: The cross-sections for electroweak-ino production at the LHC. We have included both LSP pair production and, in the case of wino and Higgsino LSP, the production of closely degenerate states, as a function of M_χ. We choose the squark mass to be 750 GeV. Right: the fraction of events with 1, 2, 3 and 4 lepton jets within the central region $|\eta| < 2.4$.

decay $\chi_0 \to h_{DM}\chi_{DM}$ is fairly large. The resulting event geometry is striking and is depicted schematically in Fig. 22).

In the left panel of Fig. 23, we show the rate of electroweak-ino pair production at the Tevatron. In the case of pure wino-like and Higgsino-like LSP, we have also included the production of the closely degenerate charginos and neutralinos. We then decay each LSP into a pair of lepton jets and study their kinematics. At the Tevatron, the neutralinos and charginos produced from $q\bar{q}'$ initial states are expected to have small boosts. Therefore, the majority of the resulting lepton jets are expected to be very central as shown in the right panel of Fig. 23, where we have required $|\eta| < 2.4$ for the lepton jets. In addition we see that the majority $> 90\%$ of the events have 4 lepton jets within the central region as illustrated in the left panel of Fig. 23. Since the presence of such lepton jets greatly enhances the possibility of triggering on such events and separating them from the background, we estimate a reach of about 300 GeV for pure Higgsino or wino LSP at the Tevatron. The case of pure bino is still difficult because of the suppression in rate.

We have shown a similar study for the LHC in Fig. 24. At the LHC, the $q\bar{q}'$ initial state will carry significant boost. Therefore, as can be seen in the right panel of Fig. 24, there is a significant fraction of events with 3 or less lepton jets in the central region, especially for the smaller electroweak-ino mass $M_\chi \leq 400$ GeV. On the other hand, as M_χ increases, the effect of the boost quickly decreases and the fraction of events with 4 lepton jets increases. Such lepton jets will give the LHC the amazing ability to probe bino production up to $M_{\tilde{B}} \sim 1$ TeV, and wino or Higgsino production up to 2 TeV.

6.2.3 Measuring the mass of the MSSM LSP

In the sorts of SUSY events shown in Fig. 22, it is possible to use lepton jets for a measurement
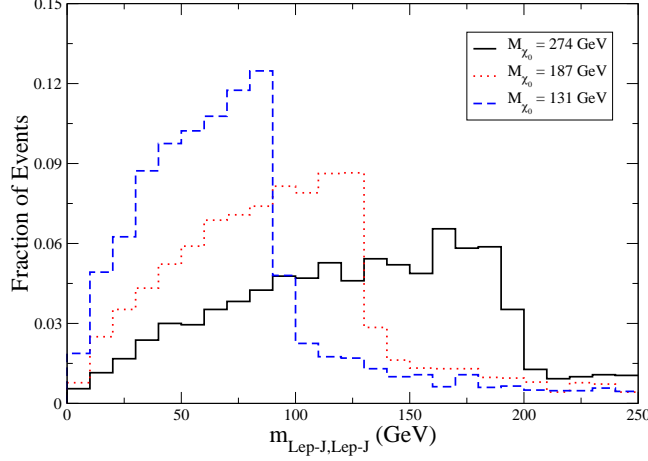
Figure 25: Forming the invariant mass of all the lepton jets in the events can lead to a measurement of the mass of χ_0, shown here as an edge at M_{χ_0}. Incorrect pairings of the lepton jets are included. However, we assumed that individual leptons are properly bunched with the correct lepton jet.

of the mass of χ_0. There are two lepton jets in each decay chain. There is a clear edge in their invariant mass distribution at M_{χ_0}, as shown in Fig. 25. This provides an absolute mass measurement and helps to remove some of the degeneracies discussed in the literature [50]. In addition, such reconstructions can be very useful in other precision measurements of the properties of the MSSM superpartners. For example, since we now have information about the LSP mass, and the direction of its decay products, it is easier to reconstruct the kinematics of the full event. In fact, we can fully recover the kinematics of the event using the same reconstruction method mentioned in the case of Z^0+jets associated production in section 6.1.5. Gaining such information will significantly improve the prospect of measuring the spin of the LSP, which can be very challenging in the conventional scenario.

7. Conclusions

In this paper we have explored numerous aspects of model building and collider phenomenology for models of dark matter in which a non-abelian dark gauge symmetry is spontaneously broken at a GeV. Assuming a minimal dark gauge group, $G_{\text{dark}} = SU(2) \times U(1)$, we have surveyed a broad class of dark Higgs sectors that break enough symmetries (charge and custodial) to be phenomenologically viable. Furthermore, in order to accommodate the XDM and/or iDM scenarios, we have included the DM as the lightest state of a gauge multiplet of G_{dark}, which we have shown can naturally acquire mass splittings of ~ 100 keV $-$ 1 MeV, from radiative corrections and higher-dimensional operators.

We have also argued that an attractive option is to include supersymmetry, since this automatically generates a GeV symmetry breaking scale for the dark sector. In particular, we have proposed a novel mechanism whereby the $U(1)$ kinetic mixing alone generates an

effective FI term for the dark hypercharge D-term, which in turn can induce SUSY breaking or a super-Higgs mechanism in the dark sector at a GeV. We emphasize that these GeV scale contributions will be always present, even if we choose to communicate SUSY breaking to the dark sector via more conventional gauge or gravity mediation. Also, in this scenario the dark matter multiplet can be uncharged under the SM gauge group, allowing us to evade astrophysical constraints which may arise if the dark matter can annihilate directly to $W^\pm$ (which in turn decay to $\bar{p}$ and π^0).

We have also given a detailed analysis of the gauge mediation scenario in which the dark matter is a $\mathbf{5} + \mathbf{\bar{5}}$ that communicates SUSY breaking from the SM to the dark sector, as originally proposed in Ref. [28]. Several benchmarks models with and without supersymmetry have been explicitly constructed here as examples.

Our model building effort is far from exhaustive. Many directions remain to be explored. For example, in the context of supersymmetric implementations, we have assumed a natural solution for the $\mu/B\mu$ problems in the dark sector. More detailed model building is certainly necessary to explore this issue further.

We have also presented a broad analysis of possible collider signatures of a non-abelian dark sector. These results are completely independent of the specific details of our benchmark models (for example how SUSY is broken, or the exact choice of dark gauge group). Indeed, our assumptions involved only the existence of a kinetic mixing and the non-abelian gauge symmetry of the dark sector.

We have found that dark gauge bosons can be produced via kinetic mixing with the photon in processes analogous to prompt photon production in the SM (with an additional suppression of ϵ^2 on the rate). At the same time, this kinetic mixing also implies that dark gauge bosons can be produced by rare Z decays with a branching ratio of the order ϵ^2. The dark sector states dominantly decay into SM $e^\pm$ and $\mu^\pm$. With GeV invariant masses and typically large boosts of $\sim$ 10s GeV, these leptons typically form highly collimated "lepton jets". Indeed, a key feature of a non-abelian dark sector is that it produces lepton jets with multiple leptons, easily > 2 and possibly even 8. We expect that identifying more than 2 leptons in a given lepton jet will significantly suppress backgrounds and enhance signal observability. With this in mind, there is the promising possibility that such signals may be observed at the LHC or possibly even at the Tevatron. Prompt photon-like production has a larger rate than rare Z-decay, but the latter is a complementary process that provides more handles on the dark sector. The 2-body decays of the dark sector into SM leptons do not generally lead to displaced vertices. However, such displaced vertices are in general possible in the case of a 3-body decay. If the 3-body decay is further suppressed by mass splitting or if the lightest scalar is the lightest dark state, then such a decay leads to $\not{E}_T$ which is collimated with the lepton jet. Several such displaced vertices will produce uncorrelated impact parameters of decay products.

In supersymmetric implementations, the lightest superpartner (LSP) of the MSSM will decay further into the dark sector states which inevitably leads to lepton jets. This offers the possibility of detecting direct electroweak gaugino production with enhanced reach both

at the Tevatron and at the LHC. We can also take advantage of the unique kinematics of such SUSY cascades to perform absolute mass measurements of the MSSM's LSP. Missing energy is in general present in these events even if the gravitino is lighter than the lightest dark superpartner (LDSP), since the decay of the LDSP$\to \tilde{G} + X$ has a lifetime much longer than the detector timescale.

We would like to emphasize that we have only outlined the leading features of the collider phenomenology of this scenario. More detailed simulations are needed to obtain fully accurate estimates of the SM background with multiple collimated leptons. There are of course variations of the models considered here which could have new types of signals. An R-parity violating scenario could have more intriguing decay chains into the dark sector [28]. The small coupling between the observable sector and the dark sector could come with other types of "portals", such as the ones presented in Refs. [1, 51], which could lead to new collider signals. We observed that the modification of the SUSY decay chain will in general lead to better kinematical reconstruction. Such new information could result in precision measurements of other properties of the MSSM superparters involved in the decay chain. New techniques for taking advantage of such new information are certainly worthy of development.

A. Kinetic Mixing

In this appendix we give a detailed description of kinetic mixing and its effect on dark sector/SM couplings. To begin, we consider the non-SUSY case. Following the proposal of [1], we couple the dark sector to the SM via a gauge kinetic mixing between the dark and SM hypercharges (see Eq. (2.1)), much like what happens in the SM between the photon and the rho meson. This scheme is attractive because it does not break any symmetries of the SM and is hence less phenomenologically constrained. Moreover, since this operator is marginal, it can be generated at a very high scale, and will persist in the infrared. This implies that if both $U(1)$'s are fundamental, the kinetic mixing is a UV boundary condition sensitive to physics at the highest scales. But if either $U(1)$ is ultimately embedded in a GUT, kinetic mixing is only induced below the GUT scale by fields charged under both $U(1)$'s. For example, by integrating out a multiplet of heavy fields Φ_i of mass M_i that is charged under both dark and SM hypercharge, we find that

$$\epsilon = -\frac{g_Y g_y}{16\pi^2} \sum_i Q_i q_i \log \left(\frac{M_i^2}{\mu^2} \right) \tag{A-1}$$

where Q_i and q_i are the charges of Φ_i under dark and SM hypercharge, and μ is the renormalization scale. If for example, SM hypercharge is generated by symmetry breaking of some GUT group under which Φ_i is charged, then $\sum_i Q_i = 0$. If this multiplet has uniform q_i charge, then the μ dependence cancels and the argument of the log becomes some ratio of scales in the multiplet, M/M'. For reasonable sizes of g_Y and g_y, and a log contribution $\log M/M' \sim 1$, this implies $\epsilon \sim 10^{-4} - 10^{-3}$, which is in the right range to explain DAMA.

Next, let us consider how the kinetic mixing induces couplings between the dark sector and SM. At the electroweak scale, the terms involving the kinetic mixing are

$$\mathcal{L}_{\text{gauge mix}} = -\frac{1}{4}W_{3\mu\nu}W_3^{\mu\nu} - \frac{1}{4}B_{\mu\nu}B^{\mu\nu} - \frac{1}{4}b_{\mu\nu}b^{\mu\nu} + \frac{\epsilon}{2}B_{\mu\nu}b^{\mu\nu} \tag{A-2}$$

$$= -\frac{1}{4}Z_{\mu\nu}Z^{\mu\nu} - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} - \frac{1}{4}b_{\mu\nu}b^{\mu\nu} + \frac{\epsilon}{2}(\cos\theta_W F_{\mu\nu} - \sin\theta_W Z_{\mu\nu})b^{\mu\nu} \tag{A-3}$$

where $F_{\mu\nu}$ and $Z_{\mu\nu}$ are the fields strengths for the SM photon and Z boson, and in the second line we have gone from gauge eigenstate to mass eigenstate. Performing a field redefinition on the photon and the dark hypercharge gauge boson

$$A_\mu{}' = A_\mu - \epsilon\cos\theta_W b_\mu \tag{A-4}$$

$$b_\mu{}' = b_\mu + \epsilon\sin\theta_W Z_\mu \tag{A-5}$$

removes the kinetic mixing between the photon and Z, and removes the kinetic mixing between the b and Z up to order ϵ^3. In addition, these shifts will modify the gauge-current couplings, $A_\mu J_{\text{em}}^\mu + Z_\mu J_Z^\mu + b_\mu J_b^\mu + w_\mu J_w^\mu$, as well as the gauge boson mass matrix. Since the photon is exactly massless, the shift of A has no effect on the mass matrix and simply couples b to the electromagnetic current of the SM. This is precisely the channel that will generate the leptons seen in astrophysical data.

Analogously, the shift of b induces a coupling of the Z boson to the dark sector b current. However, unlike the photon, the b actually acquires a mass at $\sim$ GeV, and furthermore mixes maximally with the dark w's. For this reason shifting b induces a new mass mixing term between all the dark gauge bosons and Z_μ of order $\epsilon m_b^2/m_Z^2 = (1\text{ GeV}/100\text{ GeV})^2 \times \epsilon$. This in turn generates a mass suppressed coupling between b and the Z current and between Z and the w current! Thus, after removing the kinetic mixing, all the terms that couple the SM to the dark sector are

$$\mathcal{L}_{\text{coupling}} = \epsilon b_\mu \left(\cos\theta_W J_{\text{em}}^\mu + \mathcal{O}(m_b^2/m_Z^2)J_Z^\mu\right) + \epsilon Z_\mu \left(-\sin\theta_W J_b^\mu + \mathcal{O}(m_b^2/m_Z^2)J_w^\mu\right) \tag{A-6}$$

where we have suppressed the mixing angles corresponding to the higher order contributions.

If we now add SUSY, then the kinetic mixing becomes the expression shown in Eq. (4.1). This induces a mixing term for gauginos and D-terms. Since we have already considered the D-term mixing in Section 4.1, we focus on the gauginos. The new term is

$$\mathcal{L}_{\text{gaugino mix}} = -2i\epsilon\lambda_{\tilde{b}}^\dagger \bar{\sigma}^\mu \partial_\mu \lambda_{\tilde{B}} + \text{h.c.} \tag{A-7}$$

38

where $\lambda_{\tilde{b}}$ and $\lambda_{\tilde{B}}$ are the dark and MSSM bino, respectively. Once again, since the dark bino is effectively massless at the electroweak scale, it is natural to shift it by $\lambda_{\tilde{b}} \to \lambda_{\tilde{b}} + \epsilon \lambda_{\tilde{B}}$. This yields a coupling term

$$\mathcal{L}_{\text{coupling}} = \epsilon \left(\lambda_{\tilde{B}} \tilde{J}_b + \mathcal{O}(M_{\tilde{b}}/M_{\tilde{B}}) \lambda_{\tilde{b}} \tilde{J}_B \right) \tag{A-8}$$

$$\tilde{J}_b = g_y \sum_i q_i \tilde{h}_i^\dagger h_i \tag{A-9}$$

$$\tilde{J}_B = g_Y \sum_i Q_i \tilde{H}_i^\dagger H_i \tag{A-10}$$

where $\tilde{J}_b$ and $\tilde{J}_B$ are the fermionic components of the dark and SM hypercharge supercurrents. Here the term that is $\mathcal{O}(M_{\tilde{b}}/M_{\tilde{B}})$ arises from new mass mixing terms that arise from the gaugino shift.

B. Conditions for Charge Breaking

In general it is straightforward to achieve "electroweak" breaking for $G_{\text{dark}} = SU(2) \times U(1)$, simply by introducing a negative mass squared at the origin of Higgs field space. In the non-SUSY case this tachyon is inserted by hand, while in the SUSY case it arises naturally in the gauge mediation or kinetic mixing mediation scenarios mentioned in this paper.

However, breaking G_{dark} completely, i.e. breaking charge, is a more difficult task. To see this, let us first consider the two Higgs doublet model. We can parameterize the vevs by

$$h_1 = v_1 \begin{pmatrix} \cos\alpha \\ \sin\alpha \end{pmatrix}, \qquad h_2 = v_2 \begin{pmatrix} 0 \\ 1 \end{pmatrix} \tag{B-1}$$

where h_1 and h_2 are $\mathbf{2}_{-1/2}$ and $\mathbf{2}_{1/2}$, respectively. For simplicity we have assumed that CP is preserved, and we have applied an $SU(2)$ transformation to rotate h_2 into a single real component. For a given Higgs potential it is possible to determine whether charge is broken by considering the effective potential for the charge breaking angle α. Charge is preserved only if $\alpha = 0$ or π at the minimum of the potential. Since $|h_1|^2$ and $|h_2|^2$ are independent of α, only two renormalizable potential terms can contribute: $|h_1^T \epsilon h_2|^2$ and $h_1^T \epsilon h_2$. Naively, $|h_1^\dagger h_2|^2$ contributes as well, but this term can be written as $|h_1^\dagger h_2|^2 = |h_1|^2 |h_2|^2 - |h_1^T \epsilon h_2|^2$. Expanding these contributions in terms of α, the effective potential becomes

$$V_{\text{eff}}(\alpha) = -\frac{1}{2} A \cos^2\alpha + B \cos\alpha \tag{B-2}$$

where A and B are a function of $v_{1,2}$ and the couplings. There is an extremum at $\alpha = \arccos B/A$. Checking that this point is stable, we find that a necessary condition for charge breaking is $A < 0$ and $|B/A| < 1$.

Next, let us consider this in an example. In the MSSM, the quartic couplings are fixed by the D-terms, which turns out to fix $A = g^2 > 0$, which is why charge is left unbroken.

However, it is possible to push A below zero by introducing appropriate quartic contributions to the MSSM. In our SUSY benchmarks we accomplish this by including triplets. The one triplet SUSY benchmark has the superpotential:

$$W = \mu_\Phi \text{Tr} \left(\Phi \Phi \right) + \mu H_1^T \epsilon H_2 + \lambda H_1^T \epsilon \Phi H_2 \tag{B-3}$$

Let us consider the case where the triplet is heavy and we can integrate it out; this yields an effective theory of doublets in which charge breaking is simply determined. Our SUSY benchmarks are not in this decoupling limit, but nonetheless the physics of charge breaking in the low-energy doublet model appears to persist even as the triplet mass is lowered. Integrating out the triplet yields a quartic for the doublet with a coupling of $\lambda^2 m_\Phi^2 / \mu_\Phi^2$, where m_Φ is the soft mass for the triplet. Thus in $V_{\text{eff}}(\alpha)$ for this model we find

$$A = g^2 + \lambda^2 \frac{m_\Phi^2}{\mu_\Phi^2} \tag{B-4}$$

We see that $A < 0$ only if there is a negative soft mass for the triplet that is appropriately large. While this can be easily engineered using gauge mediation, this not always possible generically. For example, if SUSY breaking is communicated in this theory via kinetic mixing mediation, then m_Φ^2 is not generated at the leading order, since Φ is a singlet under the dark hypercharge. On the other hand, we can easily remedy this by charging Φ under dark hypercharge—however, for anomaly cancellation we must also introduce a second triplet of opposite charge. In this model of two complex triplets, the kinetic mixing mediation will generate soft masses for the triplets.

C. $\mathbb{Z}_2$ Symmetry in the $\tan \beta = 1$ Limit

When $\tan \beta = 1$, there is an enhanced $\mathbb{Z}_2$ symmetry of the two Higgs doublet model that makes it phenomenologically inviable. In particular, under this symmetry the $w_{\pm \mu}$ have charge -1 and the z_μ and a_μ (dark photon) have charge $+1$. Since only z and a contain a component of the dark hypercharge, b, only these states couple to SM electric charge. On the other hand, transitions between the different dark matter states are mediated by $w_\pm$ alone. Thus it is necessary to break this $\mathbb{Z}_2$.

The origin of this $\mathbb{Z}_2$ is as follows. When $\tan \beta = 1$, then h_1 and h_2 have the same magnitude; thus, they can be simply thought of as two spinors that correspond to two unit vectors in 3-space. Next, h_1 and h_2 uniquely define a third direction which bisects the angle α between them. Rotations of 180° around this axis, followed by $h_1 \leftrightarrow h_2$, leave the vacuum invariant. Thus, all states in the low-energy theory are eigenstates of this $\mathbb{Z}_2$. Since this $\mathbb{Z}_2$ is a subgroup of $SU(2)$, it acts nicely on $\{w_1, w_2, w_3\}$. It is obvious by choosing a basis where w_3 points along the axis of rotation, that two of the $SU(2)$ gauge bosons are odd under this $\mathbb{Z}_2$ and the remaining one is even. Thus, the latter is the only state that can mix with the dark hypercharge, b, since it is also $\mathbb{Z}_2$ neutral.

D. Supersymmetric Contribution to Mass Splitting

In this appendix we present the supersymmetric contributions to the mass corrections of a Dirac fermion, Ψ. We show that in the parameter region we are interested in, those are negligible. These results are well-known (see for example [34]) and are presented here for completeness. We begin with the non-supersymmetric contributions. Consider, therefore, a theory with a SM-like weak gauge-group $SU(2) \times U(1)$ which in general is broken down to nothing at some scale. Also, for simplicity, we take the Dirac fermion to be charged as $\mathbf{2}_{1/2}$ with a mass M_Ψ much larger than the Higgs scale of the theory. Similar conclusions hold for any representation of the gauge-group. At low energies, the masses of the two components are split. There are both wave-functions and mass insertion diagrams, given by,

$$= -\frac{ig_A^2}{8\pi^2}\slashed{p}\int dx\, x \log \Delta \qquad \text{(D-1)}$$

$$= \frac{ig_A^2}{4\pi^2}M_\Psi \int dx \log \Delta \qquad \text{(D-2)}$$

where $\Delta = \left((1-x)^2 M_\Psi^2 + x M_A^2\right)$, M_A is the mass of the gauge boson, and g_A is its coupling to the fermion. We neglected the divergent part since it cancels when considering the mass splitting. The propagating gauge boson is any one of the four massive vector bosons.

For the rest of this section we consider the simple case where one gauge boson is left massless. In that case we have,

$$\delta M_\Psi = \frac{\alpha M_\Psi}{2\pi}\int dx \log\left(1 + \frac{x M_Z^2}{(1-x)^2 M_\Psi^2}\right) \xrightarrow{M_\Psi \gg M_Z} \frac{\alpha M_Z}{2} \qquad \text{(D-3)}$$

Where α is associated with the massless gauge boson coupling to matter. For a general multiplet of the gauge group the mass splitting between two eigenstates, i and j of T_3 is given by,
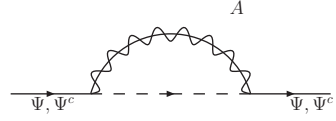
$$\delta M_{ij} = \frac{\alpha}{2}(q_i^2 - q_j^2)M_z \qquad \text{(D-4)}$$
$$- \frac{\alpha_2}{2}\left((T_i^3)^2 - (T_j^3)^2\right)(M_z - M_w),$$

where the notation is explained after Eq. (3.3) in Section 3. When the "photon" is also massive the correction goes as the splitting between the gauge boson masses,

$$\delta M_\Psi \approx \frac{\alpha}{2}(M_Z - M_\gamma) + \dots \qquad \text{(D-5)}$$

However, the precise formula requires the vector boson mass eigenstates and is not simple in general.

The supersymmetric contribution is through a similar loop to the wave-function renormalization above only with a gaugino - slepton loop replacing the gauge boson - lepton propagators[15],

 $$= -\frac{ie^2}{8\pi^2}\not{p}\int dx(1-x)\log\tilde{\Delta} \qquad \text{(D-6)}$$

with $\tilde{\Delta} = \left((1-x)^2 M_{\tilde{\Psi}}^2 + x M_{\tilde{A}}^2\right)$. The contribution to the splitting is then,

$$\delta M_\Psi = \frac{\alpha M_\Psi}{2\pi}\int dx \log\left(\frac{(1-x)^2 M_{\tilde{\Psi}^+}^2 + x M_{\tilde{W}}^2}{(1-x)^2 M_{\tilde{\Psi}^0}^2 + x M_{\tilde{\gamma}}^2}\right) \qquad \text{(D-7)}$$

where we used $M_{\tilde{W}}$, $M_{\tilde{\gamma}}$ ($M_{\tilde{\Psi}^+}$, $M_{\tilde{\Psi}^0}$) casually to designate the charged and neutral gauge bosons (leptons). Clearly, in the limit where all the masses are equal the integral vanishes. Therefore, the only contribution to the splitting comes from possible differences between $M_{\tilde{W}}^2$ and $M_{\tilde{\gamma}}^2$ (or $M_{\tilde{\Psi}^+}$ and $M_{\tilde{\Psi}^0}$). If we denote the soft supersymmetry contribution to the gaugino masses by M_λ we have,

$$M_{\tilde{W}}^2 - M_{\tilde{\gamma}}^2 \approx 2 M_\lambda M_W \qquad \text{(D-8)}$$

Therefore, this contribution to the mass splitting in Eq. (D-7) is suppressed by M_λ/M_Ψ compared with the non-supersymmetric contribution in Eq. (D-5).

Another possibility is that the charged slepton is split from the neutral one by $SU(2)$ D-term contributions. However, those contributions are to the mass squared. Writing $M_{\tilde{\Psi}^+}^2 - M_{\tilde{\Psi}^0}^2 \approx M_W^2$ we see that again the contribution to the mass splitting is suppressed by M_W/M_Ψ with respect to Eq. (D-5).

References

[1] N. Arkani-Hamed, D. P. Finkbeiner, T. Slatyer, and N. Weiner, *A Theory of Dark Matter*, 0810.0713.

[2] J. Chang *et al.*, *An excess of cosmic ray electrons at energies of 300.800 GeV*, Nature **456** (2008) 362–365.

[3] O. Adriani *et al.*, *Observation of an anomalous positron abundance in the cosmic radiation*, 0810.4995.

[4] O. Adriani *et al.*, *A new measurement of the antiproton-to-proton flux ratio up to 100 GeV in the cosmic radiation*, 0810.4994.

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[15]This way of organizing the diagrams makes it clear that even in the supersymmetric limit, the mass insertion diagram is not cancelled against anything else and the mass splitting is physical. This may appear to be in conflict with known renormalization theorems of the superpotential. However, it is important to note that such theorems are not manifest in the Wess-Zumino gauge which is used to compute the splitting. The splitting is a physical effect, but the precise diagrams involved in supersymmetry is a matter of gauge choice. For a more detailed discussion of the issues, see Ref. [52]

[5] D. P. Finkbeiner, *Microwave ISM Emission Observed by WMAP*, *Astrophys. J.* **614** (2004) 186–193, [`astro-ph/0311547`].

[6] D. P. Finkbeiner, *WMAP microwave emission interpreted as dark matter annihilation in the inner Galaxy*, `astro-ph/0409027`.

[7] G. Dobler and D. P. Finkbeiner, *Extended Anomalous Foreground Emission in the WMAP 3-Year Data*, *Astrophys. J.* **680** (2008) 1222–1234, [`0712.1038`].

[8] I. Cholis, G. Dobler, D. P. Finkbeiner, L. Goodenough, and N. Weiner, *The Case for a 700+ GeV WIMP: Cosmic Ray Spectra from ATIC and PAMELA*, `0811.3641`.

[9] E. A. Baltz *et al.*, *Pre-launch estimates for GLAST sensitivity to Dark Matter annihilation signals*, *JCAP* **0807** (2008) 013, [`0806.2911`].

[10] B. Holdom, *Two U(1)'s and Epsilon Charge Shifts*, *Phys. Lett.* **B166** (1986) 196.

[11] M. Pospelov, *Secluded U(1) below the weak scale*, `0811.1030`.

[12] M. Pospelov, A. Ritz, and M. B. Voloshin, *Secluded WIMP Dark Matter*, *Phys. Lett.* **B662** (2008) 53–61, [`0711.4866`].

[13] J. Hisano, S. Matsumoto, M. Nagai, O. Saito, and M. Senami, *Non-perturbative Effect on Thermal Relic Abundance of Dark Matter*, *Phys. Lett.* **B646** (2007) 34–38, [`hep-ph/0610249`].

[14] J. Hisano, S. Matsumoto, M. M. Nojiri, and O. Saito, *Non-perturbative effect on dark matter annihilation and gamma ray signature from galactic center*, *Phys. Rev.* **D71** (2005) 063528, [`hep-ph/0412403`].

[15] M. Cirelli, A. Strumia, and M. Tamburini, *Cosmology and Astrophysics of Minimal Dark Matter*, *Nucl. Phys.* **B787** (2007) 152–175, [`0706.4071`].

[16] J. March-Russell, S. M. West, D. Cumberbatch, and D. Hooper, *Heavy Dark Matter Through the Higgs Portal*, *JHEP* **07** (2008) 058, [`0801.3440`].

[17] M. Pospelov and A. Ritz, *Astrophysical Signatures of Secluded Dark Matter*, `0810.1502`.

[18] I. Cholis, L. Goodenough, and N. Weiner, *High Energy Positrons and the WMAP Haze from Exciting Dark Matter*, `0802.2922`.

[19] D. P. Finkbeiner and N. Weiner, *Exciting Dark Matter and the INTEGRAL/SPI 511 keV signal*, *Phys. Rev.* **D76** (2007) 083519, [`astro-ph/0702587`].

[20] D. Tucker-Smith and N. Weiner, *Inelastic dark matter*, *Phys. Rev.* **D64** (2001) 043502, [`hep-ph/0101138`].

[21] G. Weidenspointner *et al.*, *The sky distribution of positronium annihilation continuum emission measured with SPI/INTEGRAL*, `astro-ph/0601673`.

[22] **CDMS** Collaboration, D. S. Akerib *et al.*, *Limits on spin-independent WIMP nucleon interactions from the two-tower run of the Cryogenic Dark Matter Search*, *Phys. Rev. Lett.* **96** (2006) 011302, [`astro-ph/0509259`].

[23] **CDMS** Collaboration, Z. Ahmed *et al.*, *A Search for WIMPs with the First Five-Tower Data from CDMS*, `0802.3530`.

[24] **XENON** Collaboration, J. Angle *et al.*, *First Results from the XENON10 Dark Matter Experiment at the Gran Sasso National Laboratory*, *Phys. Rev. Lett.* **100** (2008) 021303, [`0706.0039`].

[25] D. N. Spergel and P. J. Steinhardt, *Observational evidence for self-interacting cold dark matter*, *Phys. Rev. Lett.* **84** (2000) 3760–3763, [`astro-ph/9909386`].

[26] R. Dave, D. N. Spergel, P. J. Steinhardt, and B. D. Wandelt, *Halo Properties in Cosmological Simulations of Self- Interacting Cold Dark Matter*, *Astrophys. J.* **547** (2001) 574–589, [`astro-ph/0006218`].

[27] L. Ackerman, M. R. Buckley, S. M. Carroll, and M. Kamionkowski, *Dark Matter and Dark Radiation*, `0810.5126`.

[28] N. Arkani-Hamed and N. Weiner, *LHC Signals for a SuperUnified Theory of Dark Matter*, `0810.0714`.

[29] S. Chang, G. D. Kribs, D. Tucker-Smith, and N. Weiner, *Inelastic Dark Matter in Light of DAMA/LIBRA*, `0807.2250`.

[30] I. F. Ginzburg and K. A. Kanishev, *Different vacua in 2HDM*, *Phys. Rev.* **D76** (2007) 095013, [`0704.3664`].

[31] M. Kamionkowski and S. Profumo, *Early Annihilation and Diffuse Backgrounds in 1/v WIMP models*, `0810.3233`.

[32] V. N. Lebedenko *et al.*, *Result from the First Science Run of the ZEPLIN-III Dark Matter Search Experiment*, `0812.1150`.

[33] G. Angloher *et al.*, *Commissioning Run of the CRESST-II Dark Matter Search*, `0809.1829`.

[34] S. D. Thomas and J. D. Wells, *Phenomenology of Massive Vectorlike Doublet Leptons*, *Phys. Rev. Lett.* **81** (1998) 34–37, [`hep-ph/9804359`].

[35] K. R. Dienes, C. F. Kolda, and J. March-Russell, *Kinetic mixing and the supersymmetric gauge hierarchy*, *Nucl. Phys.* **B492** (1997) 104–118, [`hep-ph/9610479`].

[36] P. Meade, N. Seiberg, and D. Shih, *General Gauge Mediation*, `0801.3278`.

[37] K. M. Zurek, *Multi-Component Dark Matter*, `0811.4429`.

[38] E. Poppitz and S. P. Trivedi, *Some remarks on gauge-mediated supersymmetry breaking*, *Phys. Lett.* **B401** (1997) 38–46, [`hep-ph/9703246`].

[39] Z. Komargodski and N. Seiberg, *mu and General Gauge Mediation*, `0812.3900`.

[40] M. J. Strassler and K. M. Zurek, *Echoes of a hidden valley at hadron colliders*, *Phys. Lett.* **B651** (2007) 374–379, [`hep-ph/0604261`].

[41] T. Han, Z. Si, K. M. Zurek, and M. J. Strassler, *Phenomenology of Hidden Valleys at Hadron Colliders*, *JHEP* **07** (2008) 008, [`0712.2041`].

[42] M. J. Strassler, *On the Phenomenology of Hidden Valleys with Heavy Flavor*, `0806.2385`.

[43] M. J. Strassler, *Possible effects of a hidden valley on supersymmetric phenomenology*, `hep-ph/0607160`.

[44] T. Sjostrand, S. Mrenna, and P. Skands, *PYTHIA 6.4 physics and manual*, *JHEP* **05** (2006) 026, [`hep-ph/0603175`].

[45] **CMS** Collaboration, G. L. Bayatian *et al.*, *CMS technical design report, volume II: Physics performance*, *J. Phys.* **G34** (2007) 995–1579.

[46] **CMS** Collaboration, G. L. Bayatian *et al.*, *CMS physics: Technical design report*, . CERN-LHCC-2006-001.

[47] **ATLAS** Collaboration, *ATLAS muon spectrometer: Technical design report*, . CERN-LHCC-97-22.

[48] F. Maltoni and T. Stelzer, *MadEvent: Automatic event generation with MadGraph*, *JHEP* **02** (2003) 027, [`hep-ph/0208156`].

[49] A. Arvanitaki, S. Dimopoulos, A. Pierce, S. Rajendran, and J. G. Wacker, *Stopping gluinos*, *Phys. Rev.* **D76** (2007) 055007, [`hep-ph/0506242`].

[50] N. Arkani-Hamed, G. L. Kane, J. Thaler, and L.-T. Wang, *Supersymmetry and the LHC inverse problem*, *JHEP* **08** (2006) 070, [`hep-ph/0512190`].

[51] Y. Nomura and J. Thaler, *Dark Matter through the Axion Portal*, `0810.5397`.

[52] E. Kraus and D. Stockinger, *Nonrenormalization theorems of supersymmetric QED in the Wess-Zumino gauge*, *Nucl. Phys.* **B626** (2002) 73–112, [`hep-th/0105028`].