Cosmic Signals from the Hidden Sector

Jeremy Mardon, Yasunori Nomura, and Jesse Thaler

Center for Theoretical Physics, Department of Physics, University of California, Berkeley, CA 94720
and
Theoretical Physics Group, Lawrence Berkeley National Laboratory, Berkeley, CA 94720

Abstract

Cosmologically long-lived, composite states arise as natural dark matter candidates in theories with a strongly interacting hidden sector at a scale of 10 – 100 TeV. Light axion-like states, with masses in the 1 MeV – 10 GeV range, are also generic, and can decay via Higgs couplings to light standard model particles. Such a scenario is well motivated in the context of very low energy supersymmetry breaking, where ubiquitous cosmological problems associated with the gravitino are avoided. We investigate the astrophysical and collider signatures of this scenario, assuming that dark matter decays into the axion-like states via dimension six operators, and we present an illustrative model exhibiting these features. We conclude that the recent data from PAMELA, FERMI, and H.E.S.S. points to this setup as a compelling paradigm for dark matter. This has important implications for future diffuse gamma ray measurements and collider searches.
1 Introduction and Summary

Weak scale supersymmetry is a very attractive candidate for physics beyond the standard model. It stabilizes the weak scale against potentially large radiative corrections, leads to successful gauge coupling unification, and predicts a plethora of new particles accessible at the LHC. This framework, however, also suffers from several generic cosmological problems, associated with overproduction of gravitinos or late decay of the field responsible for supersymmetry breaking \[1, 2, 3, 4\]. In many supersymmetry breaking scenarios, this requires a rather low reheating temperature after inflation, making it difficult to explain the observed baryon asymmetry of the universe.

A simple way to avoid these cosmological problems is to assume that the gravitino is very light. If the gravitino mass satisfies

\[
m_{3/2} \lesssim O(10 \text{ eV}),
\]

(1)
then the gravitino is in thermal equilibrium with the standard model down to the weak scale. The resulting gravitino abundance is small [5] and consistent with structure formation [6]. This allows for an arbitrarily high reheating temperature, and thus baryogenesis at high energies such as thermal leptogenesis [7].

It is remarkable that this simple cosmological picture is precisely the one suggested by arguably the simplest scheme for supersymmetry breaking. Suppose that supersymmetry is dynamically broken at a scale

\[ \sqrt{F} \approx O(10 - 100 \text{ TeV}), \]

(2)
giving \( m_{3/2} = F/\sqrt{3}M_{\text{Pl}} \approx O(0.1 - 10 \text{ eV}) \), where \( M_{\text{Pl}} \approx 2.4 \times 10^{18} \text{ GeV} \) is the reduced Planck scale. Then, if the sector breaking supersymmetry is charged under the standard model gauge group, gaugino and scalar masses of order \( (g^2/16\pi^2)\sqrt{F} \approx O(100 \text{ GeV} - 1 \text{ TeV}) \) can be generated through standard model gauge loops [8, 9]. The generated squark and slepton masses are flavor universal, and thus solve the supersymmetric flavor problem. The supersymmetric Higgs mass (\( \mu \) term) can also be generated through direct interactions between the Higgs fields and the supersymmetry breaking sector [10, 11] or by the vacuum expectation value of a singlet field [12].

What are the experimental signatures of this simple supersymmetry breaking scenario beyond its indirect implications on the superparticle spectrum? In this paper we advocate that the scenario may lead to distinct astrophysical and collider signatures, due to the following features of the dynamical supersymmetry breaking sector that may appear under rather generic conditions:

**Quasi-stable states:** One immediate consequence of the present framework is that dark matter cannot be the lightest supersymmetric particle, which is the very light gravitino. Dark matter, however, can arise naturally as a (quasi-)stable state in the supersymmetry breaking sector [13]. Let \( m_{\text{DM}} \) be the mass of this state. The annihilation cross section is then naturally

\[ \langle \sigma v \rangle \approx (1/8\pi)(\kappa^4/m_{\text{DM}}^2), \]

where \( \kappa \) represents typical couplings between states in the strong sector. For \( m_{\text{DM}} \approx O(10 - 100 \text{ TeV}) \), natural values for the coupling \( \kappa \approx O(3 - 10) \) give

\[ \langle \sigma v \rangle \approx (1/8\pi)(1/\text{TeV}^2), \]

which leads to the correct thermal abundance for dark matter, \( \Omega_{\text{DM}} \approx 0.2 \).

In general, a strongly interacting sector of a quantum field theory often possesses enhanced global symmetries, such as baryon number and flavor symmetries. These symmetries can lead to “stable” states if they are not broken by the strong dynamics. It is, however, quite possible that these symmetries are not respected by physics at (much) higher energies, such as at the gravitational scale. The “stable” states then become quasi-stable states, decaying through higher dimension operators on cosmological timescales.

**Light axion-like states:** A light state appears as a pseudo Nambu-Goldstone boson when an approximate global symmetry is spontaneously broken in the strongly interacting sector. Since our strong sector is supposed to break supersymmetry, a generic argument of Ref. [14]
suggests that it may possess an accidental $U(1)$ $R$ symmetry which is dynamically broken (even if supersymmetry is broken in a local minimum). This then leads to a light $R$ axion, whose mass is typically of $O(1 - 100 \text{ MeV})$ if the mass dominantly arises from a constant term in the superpotential canceling the vacuum energy of order $F^2$ \cite{15}.

More generally, it is not hard to imagine that the sector possesses enhanced approximate global symmetries. If these symmetries are spontaneously broken, light axion-like states with masses much smaller than the dynamical scale will appear. The masses of these states are then controlled by the size of explicit breaking of the corresponding symmetries.

**Couplings to the Higgs fields:** If the supersymmetry breaking sector yields an $R$ axion, it generically couples to, and thus mixes with, the Higgs fields. This is, in fact, the case even if the Higgs fields are not directly coupled to the supersymmetry breaking sector because the holomorphic Higgs mass-squared ($B\mu$ term) obtains loop contributions from gaugino masses, which are necessarily $R$-violating. For other axion-like states, mixing with the Higgs fields can arise if the Higgs fields are directly coupled to the supersymmetry breaking sector, making the corresponding symmetries Peccei-Quinn (PQ) symmetries. This is well motivated, since such couplings are often needed to generate the $\mu$ term.

It is interesting to note that all the ingredients above appear in QCD, a known strongly coupled system in nature. For the first two, one should simply think of protons, neutrons, pions, and kaons (in the appropriate limits where quark masses are small or weak interactions are superweak). Even direct Higgs couplings do exist, although the Higgs boson is much heavier than the dynamical scale of QCD, while it is much lighter than the dynamical scale considered here.

The structures described above could be manifested in various experiments. In particular, couplings between the Higgs and axion-like states provide a potential window to probe the supersymmetry breaking sector directly. Since the sector may contain dark matter as well as light states, possible signatures may appear both in astrophysical and collider physics data, and in this paper we consider the following two classes of signatures:

**Cosmic ray signals from decaying dark matter:** If dark matter is a quasi-stable state in the supersymmetry breaking sector, its decay may lead to various astrophysical signatures. Assuming the decay occurs through the light states, the final states can be mainly leptons, explaining the excess of the positron to electron ratio observed in the PAMELA experiment \cite{16}, along the lines of \cite{17, 18, 19}. The required lifetime of order $10^{26}$ sec is obtained if the decay is caused by a dimension six operator suppressed by the unification or gravitational scale \cite{20}.

There are characteristic features for this explanation of the PAMELA excess which are being tested in current observations. First, since the decay occurs through light states, it
Table 1: The symmetry structure necessary to realize the scenario presented in this paper. A light axion-like state emerges from spontaneous breaking of an $R$ or PQ (Peccei-Quinn) symmetry, which then mixes with the Higgs sector of the standard model. Composite states in the strong sector are quasi-stable because of a $B$ (“baryon number”) or $F$ (“flavor”) symmetry. The $R$ or PQ symmetry is explicitly broken by supergravity effects or by dimension five operators suppressed by $M_I \approx O(10^9 - 10^{18} \text{ GeV})$. This gives a sufficiently large mass to the axion-like state. Explicit breaking of $B$ or $F$ is due to dimension six operators suppressed by $M_* \approx O(10^{16} - 10^{18} \text{ GeV})$, leading to dark matter decay through the light states.

Table 1: The symmetry structure necessary to realize the scenario presented in this paper. A light axion-like state emerges from spontaneous breaking of an $R$ or PQ (Peccei-Quinn) symmetry, which then mixes with the Higgs sector of the standard model. Composite states in the strong sector are quasi-stable because of a $B$ (“baryon number”) or $F$ (“flavor”) symmetry. The $R$ or PQ symmetry is explicitly broken by supergravity effects or by dimension five operators suppressed by $M_I \approx O(10^9 - 10^{18} \text{ GeV})$. This gives a sufficiently large mass to the axion-like state. Explicit breaking of $B$ or $F$ is due to dimension six operators suppressed by $M_* \approx O(10^{16} - 10^{18} \text{ GeV})$, leading to dark matter decay through the light states.

typically involves (long) cascade chains. Second, the mass of dark matter is rather large, $m_{DM} \approx O(10 \text{ TeV})$, since it arises as a quasi-stable state in the supersymmetry breaking sector. These lead to a rather broad structure in the electron plus positron spectrum in the sub-TeV region after propagation to the earth. Remarkably, we find that such a structure beautifully reproduces the spectra recently reported by the FERMI [21] and H.E.S.S. [22, 23] experiments.

**Collider signals of light axion-like states:** The couplings between the Higgs and axion-like states imply that the axion-like states also couple to the standard model gauge fields at the loop level. Such couplings may also arise from contributions from the supersymmetry breaking sector. This leads to the possibility of producing the light states at the LHC, which subsequently decay into standard model fields through mixings with the Higgs fields [21]. The final states are most likely leptons, if recent cosmic ray data are explained as described above. This may provide a relatively clean signal to discover the light states.

The couplings of the Higgs to light states also raises the possibility that the Higgs boson decays into two axion-like states. If the axion-like state decays mainly into two leptons, then this leads to a four lepton final state, whose invariant mass peaks at the Higgs boson mass. If the rate is sufficiently large, this leads to a way of seeing the axion-like state (and the Higgs boson) at hadron colliders.

The symmetry structure necessary to produce the signatures described above is summarized in Table 1. While this structure is strongly motivated by the low energy supersymmetry breaking
scenario, all that is actually required is some strong dynamics at \( \approx O(10 - 100 \text{ TeV}) \) satisfying the properties given in the table. Given that QCD already has more or less all the desired ingredients, we expect that the required structure may arise naturally in wide classes of strongly interacting gauge theories, including non-supersymmetric theories. (For a non-supersymmetric theory, the symmetry leading to a light state must be a PQ symmetry. This implies that the theory must have two Higgs doublets.) The dynamical scale of \( O(10 - 100 \text{ TeV}) \) then suggests that this sector is related to the weak scale through a loop factor. This feature is automatic in low energy supersymmetry breaking theories.

A remarkable thing is that the first signature may have already been seen in the recent cosmic ray electron/positron data. The PAMELA experiment found an unexpected rise in the positron fraction in the energy range between 10 and 100 GeV, while the FERMI experiment saw an excess of the electron plus positron flux over standard diffuse cosmic ray backgrounds in the sub-TeV region. These results suggest a new source of primary electrons and positrons with a broad spectrum extending up to a few TeV. We find that these features are very well explained by dark matter in our framework: a quasi-stable state with mass of \( O(10 \text{ TeV}) \), cascading into leptons through light axion-like states with lifetime of \( O(10^{26} \text{ sec}) \). We perform a detailed analysis for the cosmic ray data and find that a wide range for the mass is allowed for the axion-like state: it can take any value between \( 2m_e \simeq 1.0 \text{ MeV} \) and \( 2m_b \simeq 8.4 \text{ GeV} \) except for a small window between \( 2m_p \simeq 1.9 \text{ GeV} \) and \( 2m_{\mu} \simeq 3.6 \text{ GeV} \).

These data, therefore, point to the setup of Table 1 as a new paradigm for dark matter, which can be beautifully realized in the framework of low energy supersymmetry breaking. Since the precise structure of the supersymmetry breaking sector is highly model dependent, one might worry that the signatures considered here may depend on many details of the supersymmetry breaking sector, which leads to large uncertainties. This is, however, not the case. As emphasized above, the existence of the signatures depends only on basic symmetry properties of the supersymmetry breaking sector, and their characteristics are determined only by a few parameters such as the mass of dark matter and the mass and decay constant of the axion-like state. While we will provide an illustrative model as a proof-of-concept, many details of the model are unimportant for the signatures. Of course, the flip side of this is that we cannot probe the detailed structure of the supersymmetry breaking sector solely by studying these signatures. We may, however, still explore some features by carefully studying cosmic ray spectra.

The signatures described here are complementary to the information we can obtain in other methods. In the framework of low energy supersymmetry breaking, the LHC will be able to measure some of the superparticle masses. This, however, may not determine, e.g., the scale of supersymmetry breaking, since the most general supersymmetry breaking sector provides little definite prediction on superparticle masses, as recently elucidated in Ref. [25]. The existence of
the very light gravitino can give specific signals, for example, those in Ref. [26]. The signatures considered here can add even more handles. In addition to indicating the specific symmetry structure of Table 1, different final states for the axion-like state decay may also be discriminated, e.g., by future measurements of the diffuse $\gamma$-ray flux at FERMI. These will provide valuable information in exploring the structure of the supersymmetry breaking sector.

The organization of the paper is as follows. In the next section, we describe our supersymmetric setup in detail. We explain that quasi-stable states with the desired lifetimes and light axion-like states with the desired masses can naturally arise. We also discuss constraints on axion-like states, and find that a wide range for the masses and decay constants are experimentally viable and can lead to leptonic decays. In Section 3, we present an example model that illustrates some of these general points. Dark matter is a stable “meson” state in the hidden sector that decays into $R$ axions with a lifetime of $O(10^{26}$ sec). In Section 4, we perform a detailed analysis of the recent cosmic ray data, and find that the results of PAMELA, FERMI, and H.E.S.S. are very well explained. We present a general analysis in the case where dark matter decays into $e^+e^−, \mu^+\mu^−, \pi^+\pi^−\pi^0$, or $\tau^+\tau^−$ either directly or through 1-step or 2-step cascades. Implications for future diffuse $\gamma$-ray measurements are also discussed. In Section 5, we briefly discuss collider signatures associated with the existence of light states. Finally, discussion and conclusions are given in Section 6, where we mention related alternative scenarios.

2 Framework

We consider a supersymmetry breaking sector which consists of fields and interactions characterized by a scale

$$\Lambda \approx O(10 - 100 \text{ TeV}).$$

The actual spectrum of this sector could span an order of magnitude or so due to its nontrivial structure. The scale of Eq. (3) is supposed to arise dynamically through some strong gauge interactions in order to explain why the supersymmetry breaking scale, and thus the weak scale, is hierarchically smaller than the Planck scale [27]. We assume the sector contains fields charged under the standard model gauge group which directly feel supersymmetry breaking. Standard model gauge loops then generate gaugino and scalar masses in the supersymmetric standard model (SSM) sector through gauge mediation [38, 9].

The supersymmetry breaking sector may also directly interact with the Higgs fields in the superpotential. A possible form for these interactions is

$$W = \lambda_u H_u O_u + \lambda_d H_d O_d,$$

or

$$W = \lambda N H_u H_d + NO_N,$$

6
where $H_{u,d}$ and $N$ are two Higgs doublet and singlet chiral superfields, respectively, and $O_{u,d,N}$ represent operators in the supersymmetry breaking sector. These interactions can generate the $\mu$ term, and thus lead to realistic electroweak symmetry breaking \cite{10,11}. A schematic picture for the current setup can be seen in Figure 1.

We assume that the entire system of Figure 1 including both the supersymmetry breaking and SSM sectors (as well as possible direct Higgs interactions), respects some approximate global symmetry under which both the supersymmetry breaking and SSM sector fields are charged. This implies that the two sectors must have an interaction that can transmit charges of the global symmetry from one sector to the other. We find that an $R$ symmetry is a natural candidate for such a symmetry, since interactions generating gaugino masses must always transmit $R$ charges. Another simple candidate is a PQ symmetry. In the presence of direct Higgs couplings to the supersymmetry breaking sector as in Eqs. (4, 5), charges of the PQ symmetry can be transmitted. Note that $R$ and PQ symmetries are global symmetries of the SSM sector only in the limit of vanishing $\mu$ term and gaugino (and holomorphic supersymmetry breaking) masses. To generate the $\mu$ term and gaugino masses, therefore, these symmetries must be spontaneously broken by the dynamics of the supersymmetry breaking sector.

The fundamental supersymmetry breaking scale of this class of theories is of order the dynamical scale

$$\sqrt{F} \approx \Lambda,$$

yielding a light gravitino, $m_{3/2} \approx \Lambda^2/M_{Pl} \approx O(0.1 - 10 \text{ eV}).$ This solves all the cosmological problems in supersymmetric theories unless the supersymmetry breaking sector introduces its own problems\footnote{Possible problems include the system being trapped in the wrong vacuum or the appearance of stable charged or colored particles.}, which we assume not to be the case. The thermal history of the universe is normal up to a very high temperature $T \gg \Lambda$, such that the relic abundances of the quasi-stable states are determined by a standard thermal freezeout calculation.
2.1 Quasi-stable states

In addition to the $R$ and PQ symmetries described above, the supersymmetry breaking and SSM sectors can have additional independent global symmetries. For example, the SSM sector has a baryon number symmetry, at least if $R$ parity is conserved (which we need not assume here). Similarly, the supersymmetry breaking sector may possess accidental global symmetries, and since this sector is assumed to feel strong gauge interactions, natural possibilities are “baryon number” and “flavor” symmetries. If the global symmetry is not spontaneously broken by the dynamics of this sector, the lightest state charged under that symmetry is stable. This state can then be dark matter, with an abundance determined by its annihilation cross section.

Since dark matter is the lightest state charged under a global symmetry, it sits at the lowest edge of the spectrum of the corresponding charged states. It is therefore natural to expect that the dark matter mass is in the range

$$m_{DM} \approx O(10 \text{ TeV}).$$

In fact, we will see later that a mass of this size reproduces well the observed electron/positron spectrum at PAMELA, FERMI, and H.E.S.S. experiments. The annihilation cross section of dark matter is naturally of order

$$\langle \sigma v \rangle \approx \frac{1}{8\pi \frac{\kappa^4}{m_{DM}^2}},$$

where $\kappa$ represents typical couplings between states in the strong sector, and we have assumed two-body annihilation. For $\kappa \approx 3$, this gives the cross section needed to reproduce the observed dark matter abundance, $\Omega_{DM} \simeq 0.2$. Such a value for $\kappa$ is quite natural for couplings between hadronic states in a strongly interacting sector.

The global symmetry ensuring the stability of dark matter is expected to be an accidental symmetry at low energies. This implies that the symmetry is not respected by physics at some higher energy $M_*$, such as the unification or Planck scale, so that the effective theory at the scale $\Lambda$ contains higher dimension operators suppressed by powers of $M_*$ that do not respect the global symmetry. This situation is precisely analogous to baryon number in the standard model embedded in grand unified theories (with $M_*$ identified with the unification scale), or strangeness in QCD (with $M_*$ identified with the weak scale). Dark matter can then decay with cosmologically observable timescales, depending on the dimension of the leading symmetry-violating operators.

Consistent with gauge coupling unification, we take $M_*$ to be around the unification or Planck scale, $M_* \approx O(10^{16} - 10^{18} \text{ GeV})$. Then, to have observable signatures in cosmic rays, the dimension of relevant symmetry violating operators must be six, as illustrated in Ref. [20] in the case where $m_{DM} \approx O(\text{TeV})$ and $M_* \approx O(10^{16} \text{ GeV})$. With dimension six decay operators, we find

$$\tau_{DM} \approx 8\pi \frac{M_*^4}{m_{DM}^5} \simeq 2 \times 10^{25} \text{ sec} \left(\frac{M_*}{10^{17} \text{ GeV}}\right)^4 \left(\frac{10 \text{ TeV}}{m_{DM}}\right)^5,$$
for two-body decays, and
\[
\tau_{\text{DM}} \approx 128\pi^3 \frac{M_*^4}{m_{\text{DM}}^6} \simeq 3 \times 10^{27} \text{ sec} \left( \frac{M_*}{10^{17} \text{ GeV}} \right)^4 \left( \frac{10 \text{ TeV}}{m_{\text{DM}}} \right)^5,
\]
(10)
for three-body decays. The required lifetime to fit the PAMELA, FERMI, and H.E.S.S. data through dark matter decay is of order $10^{26}$ sec, which is consistent with the value of $M_*$ taken here.

### 2.2 Decay of quasi-stable states

There are several possible ways for the decay of dark matter to be caused by dimension six operators. One simply needs small breaking of the global symmetries protecting dark matter, and generically decays will proceed via some kinematically-allowed but symmetry-violating channel. In particular, dark matter can decay into light axion-like states and/or the gravitino. As long as the dimension six operators are suppressed by the unification or Planck scale, they can even explicitly violate the $R$ or PQ symmetry that gives rise to the axion-like state, since such explicit breaking gives very small masses compared to the contributions considered in Section 2.3.

Suppose that interactions in the supersymmetry breaking sector are asymptotically free, so that the dimensions of the operators are determined by the canonical dimensions of the elementary fields. Suppose also, for illustrative purposes, that the supersymmetry breaking sector contains a gauge group $SU(N_c)$ ($N_c \geq 3$) with $N_f$ flavor of “quark” fields $Q^i + \bar{Q}_{\bar{i}}$ ($i, \bar{i} = 1, \ldots, N_f$), where $Q^i$ and $\bar{Q}_{\bar{i}}$ are chiral superfields in the fundamental and anti-fundamental representations of $SU(N_c)$, respectively. We also assume that the quark fields have generic masses of order $\Lambda$ (or somewhat smaller) due to some $(R$-violating) dynamics in the supersymmetry breaking sector.

The $SU(N_c)$ gauge group is supposed to confine at a scale $\approx \Lambda$. This then leads to composite “meson” fields $M_j^i \sim Q^i \bar{Q}_j$, and also “baryon” and “antibaryon” fields $B \sim Q^{N_c}$ and $\bar{B} \sim \bar{Q}^{N_c}$ for $N_f \geq N_c$. Because of the nonzero quark masses, these fields have masses of order $\Lambda$ (or somewhat smaller). Moreover, in the limit that nonrenormalizable operators vanish, the states $M_j^i$ ($i \neq \bar{j}$), $B$ and $\bar{B}$ can easily be stable (except for $B$ and $\bar{B}$ for $N_c = 3$). These states are therefore good candidates for dark matter.

The lifetimes of the quasi-stable states described above are controlled by the form of higher dimension operators. Let us begin with the case where dark matter is identified with baryons (and antibaryons). The decay of dark matter can then occur through baryon number violating operators in the superpotential. These operators are dimension six if $N_c = 5^2$

\[
W \sim \frac{1}{M_*^2} QQQQ + \frac{1}{M_*^4} \bar{Q}QQ\bar{Q}Q,
\]
(11)

\footnote{In the following equations, any one of the operators is sufficient to cause dark matter decay. In particular, physics at $M_*$ need not respect a $Q \leftrightarrow \bar{Q}$ symmetry.}
in which case the lifetime of dark matter can be of order $10^{26}$ sec as needed to reproduce the electron/positron data.

An alternative possibility is that dark matter decays through holomorphic terms in the Kähler potential $K$, which can also be viewed as superpotential terms suppressed by an extra power of $M_{Pl}$ after performing a Kähler transformation. The required lifetime is obtained for $N_c = 3$:

$$K \sim \frac{1}{M_*} QQ + \frac{1}{M_*} QQQ + \text{h.c.} \quad \Rightarrow \quad W \sim \frac{\Lambda^2}{M_* M_{Pl}} QQ + \frac{\Lambda^2}{M_* M_{Pl}} QQQ,$$

where we have used the fact that the superpotential has a constant term of order $\Lambda^2 M_{Pl}$ to cancel the cosmological constant, and assumed that the renormalizable superpotential term $W \sim QQQ$ is absent, perhaps because of an $R$ symmetry. Note that in both of the cases above, the number of “colors” $N_c$ needs to be chosen appropriately to have dimension six dark matter decay.

We now consider the case where the meson states $M^i_j$ $(i \neq j)$ are dark matter. These states exist even for $N_f < N_c$, and the longevity of their lifetime could be ensured by a vector-like $U(1)^{N_f}$ symmetry that may exist in the Lagrangian at the renormalizable level. The decay of these states can be caused by dimension six operators in the Kähler potential:

$$K \sim \frac{1}{M_*} Q^i_j Q^j_k Q^k l + \frac{1}{M_*} Q^i_j Q_j Q^j k Q^k l + \frac{1}{M^2} Q^i_j Q^j Q^j k Q^k l.$$

Possible lower dimension operators in the superpotential $W \sim (1/M_*)Q^i_j Q^k Q^k l$ can easily be absent, for example by imposing an $R$ symmetry. An attractive feature of this possibility is that the lifetime does not depend on the number of colors $N_c$, and that the number of flavors $N_f$ need not be equal to or larger than $N_c$. This setup, therefore, can naturally be accommodated in a wide variety of gauge theories. In Section 3 we present an explicit model realizing this possibility, where the quasi-stable mesons decay via Eq. (13) into $R$ axions.

Although we have not considered them in the discussion above, in general the supersymmetry breaking sector also contains states not charged under $SU(N_c)$. Dark matter decay may then occur through these states. For example, mesons may decay through operators of the form $K \sim Q^i_j \Phi^i \Phi / M^2 + \bar{Q}^i_j \bar{Q}_j \Phi^i \Phi / M^2$, where $\Phi$ represents a generic field in the supersymmetry breaking sector. The existence of singlet fields can also change the requirement on $N_c$ for the baryon dark matter case. As long as the lowest symmetry breaking operators are dimension six, however, the existence of these operators do not affect the basic argu-

---

3A long-lived hidden sector baryon with the $SU(5)$ gauge group was considered in Ref. [28] in the context that the quarks $Q^i$ are also charged under the standard model gauge group. This model, however, does not preserve the success of perturbative gauge coupling unification, since the extra matter content charged under the standard model gauge group is too large, see e.g. [29]. Here we consider that the quarks $Q^i$ are not charged under the standard model gauge group.
Below, we assume for simplicity that operators containing only quarks dominate the decay. Whether this is the case or not is determined by physics at the scale $M_*$. This also implies that direct interactions between the supersymmetry breaking and SSM sectors, such as 

$$K \sim Q^i Q^j \Phi_{SSM}^\dagger \Phi_{SSM} / M_*^2 + \bar{Q}^i \bar{Q}^j \Phi_{SSM}^\dagger \Phi_{SSM} / M_*^2,$$

are relatively suppressed. Here, $\Phi_{SSM}$ represents SSM fields. We will comment on the case where these interactions are relevant in Section 6.

2.3 Light axion-like states

Spontaneous breaking of an approximate $R$ or PQ symmetry in the supersymmetry breaking sector leads to a pseudo Nambu-Goldstone boson. The phenomenology associated with this particle is determined largely by its mass $m_a$ and decay constant $f_a$. Suppose that relevant interactions in the supersymmetry breaking sector obey naive dimensional analysis for a generic strongly coupled theory \[30\]. In this case, we find that the decay constant is naturally somewhat ($\approx 4 \pi$) smaller than the dynamical scale $\Lambda \approx O(10 – 100 \text{ TeV})$:

$$f_a \approx O(1 – 10 \text{ TeV}). \quad (14)$$

According to naive dimensional analysis, the generic size for expectation values of the lowest and highest components of a chiral superfield $\Phi$ are given by $\langle \Phi \rangle \approx \Lambda / 4 \pi$ and $\langle F_\Phi \rangle \approx \Lambda^2 / 4 \pi$, respectively, and a generic coupling constant has the size $\kappa \approx 4 \pi$. (Here we have ignored a possible $N_c$ or $N_f$ factor associated with the multiplicity of fields, but this does not affect the basic discussion.) This implies that a generic supersymmetric mass and supersymmetry breaking mass-squared splitting in the supersymmetry breaking sector are of order $M_{mess} \approx \kappa \langle \Phi \rangle \approx \Lambda$ and $F_{mess} \approx \kappa \langle F_\Phi \rangle \approx \Lambda^2$, respectively, so that the gaugino and scalar masses generated in the SSM sector are of order $(g^2 / 16 \pi^2) F_{mess} / M_{mess} \approx (g^2 / 16 \pi^2) \Lambda$, which is consistent with Eq. (3). On the other hand, the decay constant of an axion-like state scales as $f_a \approx \langle \Phi \rangle \approx \Lambda / 4 \pi$, giving Eq. (14). Note that the suppression of $f_a$ over $\Lambda$ here is precisely analogous to the fact that in QCD the pion decay constant, $f_\pi \approx O(100 \text{ MeV})$, is an order of magnitude smaller than the characteristic QCD scale, i.e. the rho meson mass $m_\rho \approx O(1 \text{ GeV})$.

The mass of a light axion-like state is determined by the size of explicit symmetry breaking. Let us first consider the state associated with an $R$ symmetry—an $R$ axion. An interesting feature of an $R$ axion is that it has an irreducible contribution to its mass from supergravity,

\[4\] With singlets, some of the accidental symmetries are not as “automatic” as the case without singlets. For example, the low energy $U(1)^{N_f}$ flavor symmetry does not exist if the quarks couple to two or more singlets with arbitrary Yukawa couplings.

\[5\] In general, it may not be true that all the couplings in the supersymmetry breaking sector are strong. In this case, the decay constant need not obey Eq. (14); for example, it can be easily of $O(100 \text{ TeV})$. In particular, if the messenger fields (fields in the supersymmetry breaking sector charged under the standard model gauge group) feel $R$-breaking effects though perturbative interactions, then $f_a$ for an $R$ axion needs to be of $O(10 \text{ TeV})$ or larger to generate sufficiently large gaugino masses.
whose size can be determined by the strength of fundamental supersymmetry breaking $F$. This is because to cancel the cosmological constant, the superpotential must have a constant piece $\langle W \rangle \approx FM_{\text{Pl}} \approx O(\Lambda^2 M_{\text{Pl}})$, which is necessarily $R$-violating. Since $\langle W \rangle \gg \Lambda^3$, the constant piece must come from a sector other than the supersymmetry breaking sector, implying that it appears as explicit breaking from the perspective of the supersymmetry breaking sector. This provides the following contribution to the $R$ axion mass \[ m_a^2 \approx 4\pi \frac{F\langle \Phi \rangle^3}{M_{\text{Pl}} f_a^2} \approx \frac{\Lambda^3}{4\pi M_{\text{Pl}}} \approx O(10 \text{ MeV})^2 \left( \frac{\Lambda}{100 \text{ TeV}} \right)^3, \tag{15} \] where $\langle \Phi \rangle$ is a generic vacuum expectation value in the supersymmetry breaking sector, and we have used $\langle \Phi \rangle \approx \Lambda/4\pi$, $f_a \approx \Lambda/4\pi$, and $F \approx \Lambda^2/4\pi$. The uncertainty of the estimate, however, is very large, so that we can easily imagine $m_a$ in the range $O(1 \text{ – } 100 \text{ MeV})$. Unless the supersymmetry breaking sector contains another explicit breaking of the $R$ symmetry, the $R$ axion mass is given by Eq. (15).

We now consider the case of a PQ symmetry, or of additional explicit breaking of an $R$ symmetry. If the symmetry is violated by dimension five operators in the supersymmetry breaking sector, then we expect

\[ m_a^2 \approx \frac{f_a^3}{M_I} \approx O(100 \text{ MeV})^2 \left( \frac{\Lambda}{100 \text{ TeV}} \right)^3 \left( \frac{10^{14} \text{ GeV}}{M_I} \right), \tag{16} \]

where we have introduced the scale $M_I$ for physics causing the explicit breaking. We imagine $M_I$ to take a value between an intermediate scale and the gravitational scale, $M_I \approx O(10^9 \text{ – } 10^{18} \text{ GeV})$, giving $m_a \approx O(1 \text{ MeV } – 10 \text{ GeV})$. While smaller masses are possible, they are in conflict with astrophysical measurements as we will see later. The origin of the intermediate scale might be associated, for example, with the Peccei-Quinn scale for the QCD axion, the scale of the constant term in the superpotential, or the $B-L$ breaking scale generating right-handed neutrino masses.\footnote{If we use $f_a \approx \Lambda$ instead of $\Lambda/4\pi$, we obtain $m_a^2 \approx O(100 \text{ MeV})^2(\Lambda/100 \text{ TeV})^3(10^{17} \text{ GeV}/M_I)$. This gives $m_a \approx O(1 \text{ MeV } – 1 \text{ GeV})$ for $M_I \approx M_\ast$.}

Alternatively, the explicit breaking may be due to a tiny dimensionless coupling in the supersymmetry breaking sector. Note that a possible contribution from the QCD anomaly, $m_a \approx m_\pi (f_\pi/f_a) \approx O(1 \text{ keV})(100 \text{ TeV}/\Lambda)$, is small.

### 2.4 Couplings of axion-like states

The couplings of the axion-like states to the SSM sector are completely determined by the symmetry structure. We are considering the case in which the Higgs bilinear $h_u h_d$ is charged under the symmetry that leads to the light state, where $h_{u,d}$ are the lowest components of $H_{u,d}$. This is
almost always true for an $R$ symmetry (unless direct couplings of the Higgs to the supersymmetry breaking sector force vanishing charges for $H_{u,d}$), and is by definition true for a PQ symmetry. In both cases, we can take a field basis in which the axion-like state $a$ is mixed into the $h_{u,d}$ fields:

$$h_u = v_u e^{\frac{i\beta}{\sqrt{2}f_a}} a,$$

$$h_d = v_d e^{\frac{i\beta}{\sqrt{2}f_a}} a,$$

where $v_{u,d} = \langle h_{u,d} \rangle$ and $\tan \beta = \langle h_u \rangle / \langle h_d \rangle$. The distribution of $a$ inside $h_{u,d}$ is determined by the condition that $a$ is orthogonal to the mode absorbed by the $Z$ boson, and we have arbitrary chosen the $O(1)$ normalization for $f_a$. The expressions of Eqs. (17, 18) completely determine the leading-order couplings of $a$ to the standard model quarks and leptons. For example, the couplings to the up quark, down quark, and electron relevant for $a$ decay are given by

$$L = -i m_u \cos^2 \beta \sqrt{2} f_a a \Psi_u \gamma_5 \Psi_u - i m_d \sin^2 \beta \sqrt{2} f_a a \Psi_d \gamma_5 \Psi_d - i m_e \sin^2 \beta \sqrt{2} f_a a \Psi_e \gamma_5 \Psi_e,$$

where $m_{u,d,e}$ are the up quark, down quark, and electron masses. The couplings to heavier generation fermions are similar.

The couplings of $a$ to the Higgs boson $h$ can be obtained by replacing $v_{u,d}$ as

$$v_u \rightarrow v_u + \frac{\cos \alpha}{\sqrt{2}} h, \quad v_d \rightarrow v_d - \frac{\sin \alpha}{\sqrt{2}} h,$$

in Eqs. (17, 18) and plugging the resulting expressions into the Higgs kinetic terms $L = |\partial_\mu h_u|^2 + |\partial_\mu h_d|^2$. Here, $\alpha$ is the Higgs mixing angle. This leads to

$$L = \frac{c_1 v}{\sqrt{2} f_a^2} h (\partial_\mu a)^2 + \frac{c_2}{4 f_a^2} h^2 (\partial_\mu a)^2,$$

where $v = \sqrt{v_u^2 + v_d^2}$, $c_1 = \sin \beta \cos \beta (\cos^3 \beta \cos \alpha - \sin^3 \beta \sin \alpha)$, and $c_2 = \cos^4 \beta \cos^2 \alpha + \sin^4 \beta \sin^2 \alpha$. In the decoupling limit, $\alpha \approx \beta - \pi/2$, this equation reduces to

$$L = \frac{v \sin^2 2\beta}{4 \sqrt{2} f_a^2} h (\partial_\mu a)^2 + \frac{\sin^2 2\beta}{16 f_a^2} h^2 (\partial_\mu a)^2.$$

The first term is responsible for Higgs decay into two axion-like states.

There are also couplings between $a$ and the standard model gauge bosons. Their precise values depend on the matter content of the entire theory. At a given energy scale $E$, effective direct interactions between $a$ and the gauge bosons are given by

$$L = \sum_A \frac{g_A^2 c_A}{32 \pi^2 \sqrt{2} f_a} a F^A_{\mu \nu} \tilde{F}^A_{\mu \nu},$$

(23)
where $A$ runs over the gauge groups accessible at the scale $E$, and $g_A$ and $c_A$ are the gauge coupling and a coefficient of order unity. The coefficient $c_A$ encodes the contributions from physics above $E$, and is determined by the symmetry properties of the fields integrated out to obtain the effective theory. At the electroweak scale or below, we generically expect $c_A \neq 0$ for color, $A = SU(3)_C$, and electromagnetism, $A = U(1)_{EM}$. The interaction with gluons ($A = SU(3)_C$) is responsible for direct production of $a$ at hadron colliders.

2.5 Constraints on axion-like states

The axion-like states for the $R$ and PQ symmetries considered here couple to the standard model fields as the DFSZ axion [31], except for possible differences in the numerical coefficients $c_A$. Their masses, however, are much heavier because of explicit symmetry breaking, so that the experimental constraints on them are quite different from those on the DFSZ axion. Here we summarize the constraints on $m_a$ and $f_a$ for the region relevant to our discussions. For previous related analyses, see Refs. [18, 32].

The constraints on an axion-like state $a$ depend strongly on its decay mode. The decay width of $a$ into two fermions is given by

$$
\Gamma(a \to ff) = \frac{n_f c_f^2 m_f^2 m_a}{16 \pi f_a^2} \left( 1 - \frac{4m_f^2}{m_a^2} \right)^{1/2},
$$

where $n_f = 1$ and 3 for leptons and quarks, respectively, and $c_f = \sin^2 \beta$ for $f = e, \mu, \tau, d, s, b$ while $c_f = \cos^2 \beta$ for $f = u, c, t$. The decay width into two photons is

$$
\Gamma(a \to \gamma\gamma) = \frac{c_\gamma^2 e^4 m_a^3}{32 \pi (16\pi^2)^2 f_a^2},
$$

where $c_\gamma$ represents the $c_A$ coefficient for $U(1)_{EM}$ at energies below $m_e$, and $e$ is the electromagnetic gauge coupling.

For $m_a < 2m_e$, $a$ decays mainly into two photons with the decay width of Eq. (25), giving

$$
c_{\tau a \to \gamma\gamma} \simeq 5.9 \times 10^9 \text{ m} \frac{1}{c_\gamma^2} \left( \frac{1 \text{ MeV}}{m_a} \right)^3 \left( \frac{f_a}{10 \text{ TeV}} \right)^2.
$$

A possible decay of an $R$ axion into two gravitinos has the width $\Gamma(a \to \tilde{G}\tilde{G}) \approx (1/16\pi) m_a^3 f_a^2$, and is thus negligible unless $m_a \lesssim (4\pi/\alpha)^3/2m_3/2$. For $2m_e \lesssim m_a \lesssim 2m_\mu$, $a$ decays dominantly into $e^+e^-$ with the decay width of Eq. (24), giving

$$
c_{\tau a \to e^+e^-} \simeq 3.8 \times 10^2 \text{ m} \frac{1}{\sin^4 \beta} \left( \frac{10 \text{ MeV}}{m_a} \right) \left( \frac{f_a}{10 \text{ TeV}} \right)^2 \left( 1 - \frac{4m_e^2}{m_a^2} \right)^{-1/2}.
$$
For $2m_\mu < a \lesssim 800$ MeV, the $a \to \mu^+\mu^-$ mode dominates with

$$c_\tau a \to \mu^+\mu^- \simeq 3.0 \times 10^{-4} m \left( \frac{300 \text{ MeV}}{m_a} \right) \left( \frac{f_a}{10 \text{ TeV}} \right)^2 \left( 1 - \frac{4m_\mu^2}{m_a^2} \right)^{-1/2}. \quad (28)$$

In this region, the $a \to \pi\pi$ mode is suppressed by CP invariance, and $a \to \pi\pi\pi$ has the width of order $(1/128\pi^3)(m_\pi^2 m_a/f_\pi f_a^2)$, which is much smaller than $\Gamma(a \to \mu^+\mu^-)$. For $m_a \gtrsim 800$ MeV, the $a \to \rho^*\pi \to \pi\pi\pi$ and $a \to \eta\pi\pi$ modes become important, but the final states still contain a significant fraction of leptons from charged pion decay, unless $m_a \gtrsim 2$ GeV where nucleon modes start dominating. For $2m_\tau < m_a < 2m_\mu$, a will decay dominantly into taus with

$$c_\tau a \to \tau^+\tau^- \simeq 6.3 \times 10^{-8} m \left( \frac{5 \text{ GeV}}{m_a} \right) \left( \frac{f_a}{10 \text{ TeV}} \right)^2 \left( 1 - \frac{4m_\tau^2}{m_a^2} \right)^{-1/2}. \quad (29)$$

The branching ratio into $c\bar{c}$ is $\simeq 3m_\mu^2/m_a^2\tan^4\beta$, which is highly suppressed for $\tan\beta \gtrsim 2$. For $m_a > 2m_\mu$ the $a \to b\bar{b}$ mode dominates.

Rare decays of mesons provide strong constraints on axion-like states. In particular, the $K^+ \to \pi^+a$ process gives significant constraints on the region of interest. The theoretical estimate for the branching ratio is $\text{Br}(K^+ \to \pi^+a) \gtrsim 1.1 \times 10^{-8}(10 \text{ TeV}/f_a)^2$ [33]. For $m_a < 2m_\mu$, the $a$ decay length is large enough that $a$ appears as an invisible particle in $K$ decay experiments. The experimental bound on the branching ratio is then $\text{Br}(K^+ \to \pi^+a) \lesssim 7.3 \times 10^{-11}$ [34]. While the theoretical estimate has large uncertainties, this gives the rough bound

$$f_a \gtrsim 100 \text{ TeV} \quad \text{for} \quad m_a < 2m_\mu. \quad (30)$$

For $2m_\mu < m_a < m_K - m_\pi$, $a$ decays quickly into $\mu^+\mu^-$, so that the relevant experimental data is $\text{Br}(K^+ \to \pi^+\mu^+\mu^-) \simeq 1 \times 10^{-7}$ [35], which is consistent with standard model expectations. Considering that the dimuon invariant mass would be peaked at $m_a$ for a decay, $\text{Br}(K^+ \to \pi^+a)$ should be somewhat smaller than this number, giving a conservative bound of

$$f_a \gtrsim \text{a few TeV} \quad \text{for} \quad 2m_\mu < m_a < m_K - m_\pi. \quad (31)$$

Radiative decay of $Y$ also provides constraints, but the bounds are typically $f_a \gtrsim O(\text{TeV})$ (for $a \to \mu^+\mu^-$ [36]) or weaker (for $a \to \tau^+\tau^-$ [37]), and do not significantly constrain the parameter region considered here.

There are constraints from beam-dump experiments. For example, the experiment of Ref. [38] excludes $f_a \lesssim (10 - 100)$ TeV for $2m_e < m_a < 2m_\mu$. None of these experiments, however, gives as strong bounds as the ones from kaon decay given above, except for possible small islands in parameter space. Constraints from reactor experiments are also similar. They are not as strong as those from kaon decay, except that the experiment of Ref. [39] excludes a region of $f_a$ somewhat above 100 TeV for $2m_e < m_a \lesssim 10$ MeV.
Astrophysics provides strong constraints on very light axion-like states [40]. A combination of the bounds from the dynamics of the sun, white dwarfs, and horizontal branch stars excludes $m_a \lesssim 300 \text{ keV}$ for the relevant region of $f_a \approx O(1 \text{ – } 100 \text{ TeV})$. Supernova 1987A could also provide potentially strong constraints. If an $a$ produced in a supernovae were to freely escape, then its mass would be excluded up to $\approx O(1 \text{ GeV})$ with the $a$ production rates corresponding to the $f_a$ values considered here. This is, however, not the case. For $f_a \approx O(1 \text{ – } 100 \text{ TeV})$, the produced $a$ is either trapped inside the supernovae or decays quickly, so that it does not carry away significant energy. In particular, for $m_a > 2m_\mu$, $a$ immediately decays into muons, which are then thermalized quickly. Supernova 1987A, therefore, does not strongly constrain the parameter space considered here, except that the region $m_a < 2m_e$ is excluded by nonobservation of $\gamma$ rays from $a$ decays [41].

A schematic diagram summarizing the bounds on $m_a$ and $f_a$ is depicted in Figure 2 together with the dominant $a$ decay modes. For $f_a \approx O(1 \text{ – } 100 \text{ TeV})$, the most natural region avoiding all the constraints are $m_a > 2m_\mu$, although $m_a < 2m_e$ is still allowed for $f_a \approx O(100 \text{ TeV})$. For the range $m_a \approx O(1 \text{ MeV – } 10 \text{ GeV})$ considered in Eqs. (15, 16), a natural $a$ decay mode will be either $\mu^+\mu^-, \tau^+\tau^-, \pi^+\pi^-\pi^0$, or hadronic (or $e^+e^-$ for $f_a \approx O(100 \text{ TeV})$). As we will see in Section 4, a decay into $\mu^+\mu^-, \tau^+\tau^-$, or $\pi^+\pi^-\pi^0$ (or $e^+e^-$) can provide a spectacular explanation for the recently observed electron/positron signals in cosmic rays.

---

Figure 2: A schematic picture for the constraints on the $m_a$-$f_a$ plane, with the shaded region corresponding to the excluded region. Note that the actual limits on $f_a$ have $O(1)$ uncertainties, as explained in the text. The dominant $a$ decay mode for a given value of $m_a$ is also depicted.
3 Illustration

In this section, we present an example model which illustrates some of the general points discussed above. The model has quasi-stable meson fields carrying nontrivial flavor, which we identify with dark matter. We find that this dark matter, in fact, decays dominantly into $R$ axions (an $R$ axion and a gravitino if the lightest meson field is a fermion) through dimension six operators. The lifetime of dark matter can therefore be naturally of $O(10^{26} \text{ sec})$. The resulting $R$ axion decays into standard model particles as shown in Section 2.4, leading to astrophysical signatures discussed below in Section 4. We also show that under reasonable assumptions, the correct thermal abundance for dark matter can be obtained through a standard freezeout calculation.

3.1 Setup in supersymmetric QCD

Ideally, one would want to construct a complete model where supersymmetry (and $R$ symmetry) is dynamically broken and successfully mediated to the SSM sector. It is, however, notoriously difficult to realize such explicit constructions [42]. Here, we will use the example of supersymmetric QCD to show that quasi-stable dark matter candidates can arise from strong dynamics with spontaneous $R$-symmetry breaking. In Section 3.4, we will show that this strong sector can be consistently coupled to supersymmetry breaking dynamics and to messenger fields, and we will comment on the parametric difference between the dark matter mass and the scale of supersymmetry breaking.

Our starting point here is supersymmetric QCD where the number of flavors $N_f$ is less than the number of colors $N_c$. In the ultraviolet, the matter content consists of quark fields $Q^i$ and $\bar{Q}^\dagger_i$ ($i, \bar{i} = 1, \cdots, N_f$). We couple these fields to a singlet field $S$ in the superpotential:

$$W = \lambda_i^i S Q^i \bar{Q}^\dagger_i.$$  \hspace{1cm} (32)

With only one singlet, these couplings can always be diagonalized by rotations of the quark fields, $\lambda_i^i = \lambda_i \delta_i^i$, and we can then easily see that this theory has an accidental $U(1)^{N_f}$ flavor symmetry, where each $U(1)$ factor corresponds to a vector-like rotation of the quark fields of a given flavor. The theory also possesses a non-anomalous $U(1)_R$ symmetry with

$$R(Q) = R(\bar{Q}) = 1 - \frac{N_c}{N_f}, \quad R(S) = \frac{2N_c}{N_f}.$$ \hspace{1cm} (33)

To realize our setup, the $R$ symmetry must be spontaneously broken. Here we treat its effect through the background expectation value of the $S$ field, which is sufficient to understand the properties of dark matter. We will eventually allow $S$ to be a dynamical field in Section 3.4.
Below the dynamical scale of QCD, which we take to be close to the scale charactering the entire supersymmetry breaking sector, the appropriate degrees of freedom to describe the dynamics are composite meson fields $M^i_j \sim Q^i Q_j$ interacting via a non-perturbative superpotential \[ W = \tilde{\lambda}_i \Delta S M_i + \lambda_M (N_c - N_f) \Lambda^3 \left( \frac{\Lambda^{N_f}}{\det M} \right)^{\frac{1}{N_c - N_f}}, \] (34)

where $\Lambda$ is the dynamical scale, and $M_i \equiv M^i_i$ ($i = \bar{i}$) are the diagonal mesons. The coefficient $\lambda_M$ is an unknown factor coming from canonically normalizing the meson fields, and a similar factor for the first term is absorbed into the definition of $\tilde{\lambda}_i$. Using naive dimensional analysis, the size of these coefficients are $\lambda_M \approx O((4\pi)^{-2(N_c - N_f)/(N_c - N_f)})$ and $\tilde{\lambda}_i/\lambda_i \approx O(1/4\pi)$. Note that there are no low energy baryon fields when $N_f < N_c$.

Setting $\langle S \rangle \neq 0$ but $\langle F_S \rangle = 0$, the superpotential of Eq. (34) has a stable supersymmetry-preserving, $R$-violating minimum. Defining $m_i = \tilde{\lambda}_i \langle S \rangle$ and $m = \prod_i m_i$, the minimum of the potential is at

\[ \langle M_i \rangle = \frac{\alpha_M \Lambda^2}{m_i}, \quad \langle M_j \rangle = 0 \quad (i \neq j), \] (35)

where for convenience we have defined $\alpha_M = \lambda_M (\det m/\Lambda_M^{N_f} \Lambda^{N_f})^{1/N_c}$. According to naive dimensional analysis, $m_i \approx O(\lambda \langle S \rangle / 4\pi)$ and $\alpha_M \approx O((\lambda \langle S \rangle / \Lambda)^{N_f/N_c}/16\pi^2)$, where $\lambda$ represents the size of the original couplings in Eq. (32). Since $\langle M_j \rangle = 0$ for $i \neq j$, the $U(1)^{N_f}$ flavor symmetry is unbroken, making the $M^i_j$ mesons with $i \neq j$ (for now, absolutely) stable. Therefore, the lightest components of these fields, either scalars or fermions, are dark matter candidates.

The diagonal mesons $M_i$ are unstable. Since the $R$ symmetry is assumed to be spontaneously broken, the theory has a light $R$ axion with the decay constant $f_a$ of order $\Lambda$. The couplings of the mesons to the $R$ axion are determined by symmetries, and by doing appropriate field redefinitions, these couplings can be read off from the meson kinetic terms. For the scalar components, the kinetic term is $\mathcal{L} = |\partial_{\mu} M_{ii} + ic_M (\langle M_i \rangle + M_{ii})(\partial_{\mu} a)/f_a|^2$, where $M_{ii} \equiv M_i - \langle M_i \rangle$ and $c_M$ is an $O(1)$ coefficient that depends on $R$ charges of the fields. Assuming real $\langle M_i \rangle$, this gives a coupling of $\text{Re}(M_{ii})$ to two $R$ axions, and coupling between $\text{Re}(M_{ii})$, $\text{Im}(M_{ii})$, and an $R$ axion. Therefore, $\text{Re}(M_{ii})$ decays promptly into two $R$ axions, while $\text{Im}(M_{ii})$ decays (somewhat less) promptly into three $R$ axions via an off-shell $\text{Re}(M_{ii})$. After supersymmetry breaking, the fermionic component of $M_{ii}$ has a coupling to the corresponding scalar component and a gravitino, suppressed by powers of $\Lambda$. This allows the $M_{ii}$ fermion to decay promptly into a gravitino and an $R$ axion.

---

8 We discuss a potential difference between the dynamical scale $\Lambda$ appearing here and the scale that breaks supersymmetry in Section 3.4.

9 Once the $S$ field becomes dynamical, $S$ will mix with the meson fields. This can introduce extra contributions to the couplings, especially if there is no hierarchy between the scales involved. It, however, does not affect the basic conclusion here.
3.2 Thermal relic abundance

Expanding around the minimum of the potential from Eq. (35), the mass terms for the mesons are

\[ W = \frac{\Lambda}{2\alpha_M} \sum_{ij} \left( M_{ij} M_{ji} \frac{m_i m_j}{\Lambda^2} + \frac{1}{N_c - N_f} M_{ii} M_{jj} \frac{m_i m_j}{\Lambda^2} \right), \]  

where we have used the notation \( M_{ij} \equiv (M_i^j - \langle M_i^j \rangle) \). The stable mesons (\( M_{ij} \) with \( i \neq j \)) do not have mixing terms because of the unbroken \( U(1)^{N_f} \) symmetry, while the unstable mesons (\( M_{ii} \)) have a mixing term that is suppressed by \( 1/(N_c - N_f) \). For simplicity of discussion, we will focus on the large \( N_c \) limit where there is a simple relationship between the stable and unstable meson masses:

\[ \text{mass}(M_{ij}) \simeq \sqrt{\text{mass}(M_{ii}) \text{mass}(M_{jj})}. \]  

(37)

Up to \( 1/(N_c - N_f) \) corrections, the leading interaction term for the mesons is

\[ W = -\frac{1}{3\alpha_M^2} \sum_{ijk} M_{ij} M_{jk} M_{ki} \frac{m_i m_j m_k}{\Lambda^3} + \frac{1}{4\alpha_M^3} \sum_{ijkl} M_{ij} M_{jk} M_{kl} M_{li} \frac{m_i m_j m_k m_l}{\Lambda^5}. \]  

(38)

The superpotential coupling in Eq. (38) allows for various annihilation diagrams for stable \( M_{ij} \) mesons. By the mass relation in Eq. (37), either the \( ii \)-type or \( jj \)-type mesons will be lighter than the \( ij \)-type mesons. Assuming \( m_i < m_j \), the annihilation cross section for \( M_{ij} M_{ji} \rightarrow M_{ii} M_{ii} \) scales like

\[ \langle \sigma v \rangle \sim \frac{\kappa^4}{8\pi m_{DM}^2}, \quad m_{DM} \sim \frac{1}{\alpha_M} \frac{m_i m_j}{\Lambda}, \quad \kappa \sim \frac{1}{\alpha_M^2} \frac{m_i^2 m_j}{\Lambda^3}. \]  

(39)

Using naive dimensional analysis, these quantities can be estimated as

\[ m_{DM} \sim \Lambda \left( \frac{\lambda \langle S \rangle}{\Lambda} \right)^{\frac{2-N_f}{N_c}}, \quad \kappa \sim 4\pi \left( \frac{\lambda \langle S \rangle}{\Lambda} \right)^{\frac{2N_f-2}{N_c}}, \]  

(40)

so that the required annihilation rate can naturally be obtained. For example, it is easy to obtain \( m_{DM} \sim O(10 \text{ TeV}) \) and \( \kappa \) a factor of a few, which leads to \( \langle \sigma v \rangle \sim (1/8\pi)(1/\text{TeV}^2) \), and thus \( \Omega_{DM} \simeq 0.2 \). Note that here we have not included possible multiplicity factors in naive dimensional analysis, which will somewhat decrease the value of \( \kappa \). We therefore do not prefer very large \( N_c \) in practice. Note also that, strictly speaking, our analysis is not theoretically very well under control for \( \lambda \langle S \rangle \gtrsim \Lambda \), although we expect that the basic dynamics is still as presented for \( \lambda \langle S \rangle \sim \Lambda \).

There is subtlety when annihilation occurs into states of comparable mass. If \( m_i \simeq m_j \), then the \( ij \)-, \( ii \)-, \( jj \)-type are nearly degenerate for relatively large \( N_c \). In this case, one must use the methods of Ref. [44] to properly calculate the thermally averaged annihilation rate. One can, however, still obtain the right thermal relic abundance with somewhat stronger couplings. If \( N_f \) is large, then there is a potential for multiple dark matter components having comparable abundances. The freezeout calculation can then be affected by various co-annihilation channels.
3.3 Dark matter decays

The $U(1)^N_f$ flavor symmetry which ensures the stability of dark matter arises as an accidental symmetry of the renormalizable interactions of Eq. (32). As such, it is plausible that this symmetry is not respected by physics at the unification or gravitational scale. Suppose that the leading operators encoding the high energy physics are the Kähler potential terms

$$K = \frac{1}{M_*^2} \eta_{ij}^k Q_i^j Q_i^k,$$

with arbitrary flavor structures for $\eta_{ij}^k$ in the basis where the interactions of Eq. (32) are diagonal: \( \lambda_i^j = \lambda_i \delta_i^j \). Below the scale \( \Lambda \), these operators can then be matched into

$$K = \frac{\Lambda^2}{M_*^2} c_{jk} M_i^{\dagger j} M_i^k,$$

where \( c_{jk} \) are coefficients. Using naive dimensional analysis, we find \( c_{jk} \approx O(1/16\pi^2) \) for \( \eta_{ij}^k \approx O(1) \) and \( \lambda_i \langle S \rangle \approx O(\Lambda) \). In general, one expects \( c_{jk} \) to have \( O(1) \) CP-violating phases.

The existence of the terms in Eq. (42) induces mixings between the diagonal and off-diagonal meson states, leading to the decay of dark matter. The decay width has a large suppression from the \( \Lambda^2/M_*^2 \) factor in Eq. (42), so that the lifetime of dark matter is very long. One possible decay chain arises through \( M_{ii} \rightarrow aa \), where \( a \) is the \( R \) axion. Through the mixing, this leads to

$$M_{ij} \rightarrow aa,$$

where we have assumed that dark matter \( M_{ij} \) is a scalar. For fermionic dark matter, one of the \( a \)'s must be replaced by a gravitino \( \tilde{G} \). Another possibility is that dark matter decays through the meson self interactions of Eq. (38). Again, through the mixing, this can lead to

$$M_{ij} \rightarrow M_{ii} M_{ii},$$

if it is kinematically allowed, \( \text{mass}(M_{ij}) > 2 \text{mass}(M_{ii}) \). Here, \( M_{ii} \) in the final state can be either a scalar or fermion, which subsequently decays into \( aa \) or \( a \tilde{G} \), respectively.

For the scalar dark matter decay in Eq. (43), the lifetime of dark matter can be estimated as

$$\tau_{DM} \approx \frac{8\pi f_a^4}{m_{DM}^3 \langle M \rangle^2} \left( \frac{M_*^2}{c} \right)^2 \approx 3 \times 10^{27} \text{sec} \left( \frac{1/16\pi^2}{c} \right)^2 \left( \frac{M_*}{10^{17} \text{GeV}} \right)^4 \left( \frac{10 \text{ TeV}}{\Lambda} \right)^5,$$

where \( \langle M \rangle \) is a typical meson expectation value, \( c \) represents a generic size for the coefficients \( c_{jk} \) in Eq. (42), and we have used \( m_{DM} \approx 4\pi f_a \approx \Lambda \) in the last equation. For \( M_* \) of order the

\[\text{The same argument as below applies to the operator } K = (1/M_*^2) Q_i^{k\dagger} Q_j^{l\dagger} Q_l^{k} \text{ or } (1/M_*^2) Q_i^{k} Q_j^{l\dagger} Q_l^{k}.\]

\[\text{Because of the CP-violating phases in } c_{jk}, \text{ even Im}(M_{ij}) \text{ dominantly decays into two } R \text{ axions, not three.}\]
unification scale, the lifetime is in the range required to produce observable cosmic ray signatures. While there are different parametric dependences for $\tau_{DM}$ if one uses the decay mode in Eq. (44) or if one considers fermionic dark matter decay, they all give the same estimate as Eq. (45) after using naive dimensional analysis.

3.4 Towards a more complete theory

In the above discussion, we assumed that $\langle S \rangle$ could be treated as a spurion for the spontaneous breaking of the $R$ symmetry. In a complete theory, $S$ would be a propagating degree of freedom that is part of a larger supersymmetry breaking sector. Naively coupling $S$ to the mesons as in Eq. (34), the $F_S$-term potential would change the meson vacuum structure. In particular, even if $S$ could be stabilized through terms in the Kähler potential, there is generically a mesonic runaway direction [45].

Therefore, it is important to have a proof-of-concept that a propagating $S$ field can not only obtain an $R$-violating vacuum expectation value but also couple to the mesons without introducing runaway behavior. In addition, one would like to see that appropriate messenger fields can be added to the theory to communicate supersymmetry breaking to the SSM. While a complete theory would incorporate these two effects in a completely dynamical setting with all scales set by dimensional transmutation, here we consider O’Raifeartaigh-type models in order to treat these effects modularly (with a dynamical model mentioned only briefly towards the end). We leave further model building to future work, although we emphasize that the precise details of supersymmetry breaking are largely irrelevant for the dark matter discussion.

We will also not be concerned by the specific mass hierarchies needed to realize a realistic theory. Typically, the scale of supersymmetry breaking should be closer to $O(100 \text{ TeV})$ to obtain a realistic superparticle spectrum, while the mass scale for dark matter suggested by the cosmic ray data is closer to $O(10 \text{ TeV})$. In a single scale theory with naive dimensional analysis, both mass scales are expected to coincide. However, in a complete dynamical theory of supersymmetry breaking, there may be additional structures that generate such a “little” hierarchy\textsuperscript{12}. For example, this can be realized if the $SU(N_c)$ theory in Section 3.1 is in a strongly interacting conformal window at high energies. If nontrivial dynamics kicks in at the scale $O(100 \text{ TeV})$ to break supersymmetry, the same dynamics can make the $SU(N_c)$ theory deviate from the fixed point, presumably due to decoupling of some degrees of freedom, leading to the dark matter setup considered here. The physics of $R$ breaking will be associated with the higher scale dynamics.

To see how $S$ can be made dynamical, consider adding the following interactions to Eq. (34):

$$W = \lambda_X X \left( S\bar{S} - \mu_S^2 \right) + m_S SY + \bar{m}_S \bar{S} \bar{Y}. \quad (46)$$

\textsuperscript{12}This little hierarchy may not be necessary in general if the axion-like state decays into $\pi^+\pi^-\pi^0$ or $\tau^+\tau^-$ (see Figure 4) and the number of messenger fields is relatively large.
These interactions could be the low energy description of a more complete supersymmetry breaking sector, and so we consider that \( \mu_S, m_S, \) and \( \bar{m}_S \) are roughly of order \( \Lambda \). Regardless of the \( R \) charge of the \( S \) field (i.e. regardless of \( N_f \) and \( N_c \)), there is a unique consistent \( R \) charge assignment for the new fields, with
\[
R(X) = 2, \quad R(\bar{S}) = -\frac{2N_c}{N_f}, \quad R(Y) = 2 - \frac{2N_c}{N_f}, \quad R(\bar{Y}) = 2 + \frac{2N_c}{N_f}.
\]  
(47)

With the inclusion of the \( m_S \) term, the \( F_S \) equation of motion simply sets the value of \( Y \) and does not change the potential for the mesons, so that the global minimum of the meson potential is still given by Eq. (35). The \( \lambda_X \) coupling forces \( S \) and \( \bar{S} \) to obtain vacuum expectation values, and the \( m_S \) and \( \bar{m}_S \) terms constrain \( S \) and \( \bar{S} \) from running away to infinity. As long as \( \lambda_X^2 \mu_S^2 > m_S \bar{m}_S \), this theory has a supersymmetry- and \( R \)-breaking global minimum away from the origin, with \( S \) vacuum expectation value given by
\[
\langle S \rangle = \sqrt{\frac{\bar{m}_S (\mu_S^2 - m_S \bar{m}_S)}{\lambda_X^2}}.
\]  
(48)

At tree level, the only fields that obtain \( F \)-term vacuum expectation values are \( X, Y, \) and \( \bar{Y} \), so the off-diagonal mesons only feel supersymmetry breaking at loop level. The vacuum expectation value of \( X \) is unspecified at tree level, but one expects loop effects to stabilize \( X \), though not necessarily at \( \langle X \rangle = 0 \). Since supersymmetry and \( R \) are both broken regardless of the value of \( \langle X \rangle \), this is an example of a model where supersymmetry and \( R \) are spontaneously broken at tree level.\(^{13}\)

In principle, one could calculate the full spectrum of the theory at one loop, allowing one to test whether the stable dark matter is the fermionic or bosonic component of the off-diagonal mesons. In practice, this is a nontrivial exercise given the large amount of mixing between the various new degrees of freedom. If the dynamical sector is truly strongly coupled, then loop counting would not apply, and the mesons could feel \( O(1) \) supersymmetry breaking effects depending on the details of the Kähler potential. Therefore, one would not expect a one-loop calculation to get the correct sign for the fermion/boson splitting.

We have seen that Eq. (32) can be consistently coupled to the supersymmetry- and \( R \)-breaking sector in Eq. (16). It is then straightforward to couple messenger fields to the supersymmetry breaking, using an analogous structure to Eq. (46). Regardless of how \( X \) is stabilized, the \( Y \) field obtains vacuum expectation values in both the lowest and highest components, so \( Y \) can communicate supersymmetry breaking to the SSM sector through complete \( SU(5)_{SM} \) messenger

\(^{13}\)In the language of Ref. [46], Eq. (46) is the \( \text{"}g = 0\text{"} \) superpotential which preserves a \( U(1)_R \) symmetry and spontaneously breaks a global \( U(1)_A \) symmetry. The inclusion of Eq. (32) identifies a linear combination of \( U(1)_R \) and \( U(1)_A \) as the true \( U(1)_{R'} \) symmetry, such that the \( R \) symmetry is broken.
Consider adding the superpotential

\[ W = \lambda_Y Y F \bar{F} + m_F F' \bar{F}' + \bar{m}_F \bar{F} F', \]  

(49)

where \( F \) and \( F' \) are 5s of \( SU(5)_{SM} \), and \( \bar{F} \) and \( \bar{F}' \) are \( \bar{5} \)s, and there is a consistent \( R \) charge assignment for the messengers. Taking \( m_F \bar{m}_F > \lambda_Y m_S \langle S \rangle \), we can ensure that the messengers do not develop vacuum expectation values. With this kind of messenger sector, the gaugino masses are proportional to \( F_Y^3 \) instead of \( F_Y \) [47]. They are, therefore, suppressed compared with the scalar masses for \( \lambda_Y F_Y \ll m_F \bar{m}_F \), which may not be fully desired. Moreover, one also needs to take \( m_F \simeq \bar{m}_F \) to avoid an unwanted large Fayet-Iliopoulos term for \( U(1) \) hypercharge. As a proof-of-concept, however, we see that it is possible to couple the dark matter, supersymmetry breaking, and messenger sectors together in a consistent way.

If one desires, the above model can be extended in such a way that all the dimensionful scales \( \mu_S, m_S, \bar{m}_S, m_F, \) and \( \bar{m}_F \) are generated dynamically. These scales can be replaced by a single chiral superfield \( T \), which obtains a vacuum expectation value through \( W = Z(T^2 - \mu_T^2) \). The new scale \( \mu_T \) can then be generated from dimensional transmutation, e.g., by replacing \( \mu_T^2 \) with a quark condensate as in Ref. [47]. For appropriate parameter choices, this extension leaves the O'Raifeartaigh dynamics largely intact. To connect the \( \mu_T \) scale to the \( \Lambda \) scale of supersymmetric QCD, we can introduce extra quarks \( Q' \) and \( \bar{Q}' \) of \( SU(N_c) \) that make the theory conformal at high energies. The quarks obtain masses from \( W = TQ'\bar{Q}' \), triggering the exit from a strongly-coupled fixed point. This provides an existence-proof model where all the scales are generated dynamically associated with single dimensional transmutation. The \( \mu \) term can also be generated by introducing singlet “messengers” and coupling them, together with the doublet messengers, to the SSM Higgs fields as in Ref. [10].

Note that there is a \( Z_2 \) parity in Eq. (49) under which all the messenger fields are odd. Such a parity is often present in realistic constructions for the supersymmetry breaking sector. In order to avoid the problem of unwanted colored/charged relics, however, the messenger fields (either elementary or composite) must decay. If the decay occurs through dimension five operators suppressed by \( M_* \) and if the decay products contain an SSM state, then this may explain the discrepancy of the measured \( ^7\text{Li} \) abundance from the prediction of standard big-bang nucleosynthesis, along the lines of Ref. [48]. In fact, the lowest dimension for the operators causing messenger decay can easily be five if, for example, the messenger fields are two-body bound states, such as meson states, of some strong dynamics.

4 Astrophysical Signatures

We have seen that dark matter in the present scenario naturally has the following features, which are not shared by the standard weakly interacting massive particle (WIMP) scenario:
• The mass of dark matter is of $O(10 \text{ TeV})$, which is significantly larger than the weak scale. The correct thermal abundance, however, is still obtained because of the relatively large annihilation cross section.

• Dark matter can decay through dimension six operators, and thus with lifetime of $O(10^{26} \text{ sec})$. This can lead to observable cosmic ray signatures.

• Dark matter can decay into light axion-like states (and to the gravitino if dark matter is a fermion), which in turn decays into standard model particles. The final states can naturally be only $e^+e^-, \mu^+\mu^-, \pi^+\pi^-\pi^0$, or $\tau^+\tau^-.

In fact, these features are precisely what are needed to explain recent electron/positron cosmic ray data from PAMELA, FERMI, and H.E.S.S. In this section, we demonstrate that recent astrophysical data are indeed beautifully explained in the present setup, and provide a fit to the parameters $m_{\text{DM}}$ and $\tau_{\text{DM}}$ using the observed electron/positron fluxes. We also discuss the diffuse $\gamma$-ray flux that could be seen in the near future by experiments such as FERMI.

Let us begin with a summary of the observational situation. The PAMELA experiment has recently reported an unexpected rise in the positron fraction $\Phi_{e^+}/(\Phi_{e^+}+\Phi_{e^-})$ in the energy range between about 10 and 100 GeV [16]. On the other hand, they did not see any deviation from the expected background in the antiproton data [49]. The FERMI experiment is the first to be able to measure the combined electron and positron flux with good precision and control of uncertainties over the entire range from 20 GeV to 1 TeV. Their recent data also show an excess over standard diffuse cosmic ray backgrounds [21], with a broad structure extending up to the highest energies [14]. At higher energies, data from the H.E.S.S. experiment indicates a spectral break in the combined electron and positron flux at around 1 TeV, with the spectral index increasing from $\approx 3.0$ to $\approx 4.1$ [23].

While there remain experimental uncertainties as well as difficulty in calculating astrophysical background fluxes, taken together these results suggest a new source of primary electrons and positrons with a broad spectrum extending up to a few TeV. An exciting possibility is that these are signals of annihilation or decay of galactic dark matter [52, 20], although astrophysical interpretations, e.g. in terms of nearby pulsars [53], are also possible. When interpreted in terms of dark matter, the data has certain implications:

• The structure around TeV in the combined $e^+$ and $e^-$ flux is very broad, implying that electrons/positrons do not arise directly from dark matter annihilation or decay; rather, they arise through some cascading processes.

\textsuperscript{14}The ATIC [50] and PPB-BETS [51] experiments have also reported an excess in the combined flux at around 600 GeV, although with a peaked spectral shape that does not seem to fully agree with the FERMI data. With their significant statistical and experimental uncertainties, we do not include these results in our analysis.
• The spectral cutoff around TeV, together with the broad structure, implies that the mass of dark matter is larger than TeV. In particular, if cosmic rays arise from decay of dark matter, then the mass scale is more like $O(10 \text{ TeV})$, especially if the cascade is sufficiently long.

• The absence of signal in the antiproton data implies that dark matter annihilates or decays mainly into leptons, although precisely how much nucleonic final states should be suppressed is not completely clear because of uncertainties in proton/antiproton propagation models.

These interpretations are consistent with detailed analyses performed after the recent FERMI data release [54, 55].

The fact that a naive WIMP dark matter candidate, such as the neutralino lightest supersymmetric particle (LSP), does not satisfy the above criteria is quite suggestive. In the present scenario, dark matter mass is naturally of $O(10 \text{ TeV})$, which decays with lifetime of $O(10^{26} \text{ sec})$ producing the observed $e^\pm$ signatures. The decay occurs through a light axion-like state (axion portal), so that the final states can selectively be leptons. This also ensures that the final state leptons arise through cascades, making the spectrum consistent with the latest FERMI and H.E.S.S. data. It is also worth mentioning that decaying dark matter is much less constrained than annihilating dark matter [55], which has some tension with $\gamma$-ray and neutrino observations [58, 59].

4.1 PAMELA, FERMI, and H.E.S.S. electron/positron data

There are various possible decay chains in the present scenario, depending on which state is dark matter and the decay properties of the axion-like state $a$. Here we consider two representative classes: 1-step cascade ($\phi \to aa; a \to \ell^+\ell^-$) and 2-step cascade ($\phi \to \phi'\phi'; \phi' \to aa; a \to \ell^+\ell^-$), where $\phi$ and $\phi'$ represent dark matter and another state in the supersymmetry breaking sector, respectively, and $\ell = e, \mu, \tau$. We also consider $a \to \pi^+\pi^-\pi^0$. An example with a 1-step cascade was already seen in Section 3: dark matter is (the scalar component of) the lightest “meson,” which decays into two axion-like states. The axion-like state then decays into standard model fields, as seen in Section 2.5. A 2-step cascade can easily arise if (the scalar component of) the lightest “baryon” is dark matter. After including dimension six baryon number violating operators, this state can decay into two meson states, each of which then decays into two axion-like states. Alternatively, dark matter may be a meson state which dominantly decays into other meson states, as seen in Section 3.3.

The processes just described are depicted in Figure 3. In the case where dark matter is a fermionic component, one side of some step of the cascade should be replaced by the gravitino.
Figure 3: Cascade decays of dark matter $\phi$ through an axion-like state $a$. Here, $\phi'$ is an unstable state in the supersymmetry breaking sector, and $\ell^\pm (\ell = e, \mu, \tau)$ are standard model leptons.

but this does not change the predicted signals, except that the dark matter lifetime must be rescaled by appropriate factors. Below, we present the fit of the predicted $e^\pm$ fluxes to the PAMELA and FERMI data. We find that good agreements between the predictions and data are obtained for all the cases considered, except that the 1-step $e^\pm$ case may have some tension with the H.E.S.S. data. While we only present the results for $a \rightarrow e^+e^-$, $\mu^+\mu^-$ and $\tau^+\tau^-$ here, we also performed the same analyses for $a \rightarrow \pi^+\pi^-\pi^0$. We find that the results in this case are very similar to the case of $a \rightarrow \tau^+\tau^-$, since the final state $e^\pm$ spectrum arising from charged pion decay is similar to that arising from $\tau$ decay.

Our analysis here follows that of Ref. [59], where galactic propagation of $e^\pm$ is treated by the standard diffusion-loss equation. The primary injection spectra are calculated as described there, assuming a large mass hierarchy in each cascade step. In addition, tau decays are simulated using PYTHIA 8.1 [60]. We assume an NFW halo profile [61] and use the MED propagation model parametrization given in [62]. While there is significant uncertainty in the halo profile within a few kpc of the galactic center, its effect on the predicted $e^\pm$ fluxes is small. We treat the astrophysical background as in Refs. [59, 52]: we take the parameterization of [63] and marginalize over the overall slope $P$ and normalization $A$ in the range $-0.05 < P_{e^+,e^-} < 0.05$ and $0 < A_{e^+,e^-} < \infty$. There remain substantial uncertainties in the background at the high energies explored by H.E.S.S., where primary electron fluxes depend strongly on individual sources within $\approx 1$ kpc from the Earth. We therefore do not include the H.E.S.S. data in our fit procedure.

We performed a $\chi^2$ analysis of signal plus background fluxes to the PAMELA and FERMI data. The PAMELA data at energies less than 10 GeV is strongly affected by solar modulation, and we exclude it from our analysis. The FERMI experiment released both statistical and systematic errors with their data. We conservatively combine these in quadrature, but note that this is likely an overestimation of the errors. The FERMI data is also subject to an overall systematic uncertainty in energy of $+5\%\, -10\%$, under which all energies are rescaled as $E \rightarrow rE$; we therefore also
Figure 4: Regions of best fit (at 68% C.L.) to the PAMELA and FERMI data for dark matter mass $m_{\text{DM}}$ and lifetime $\tau_{\text{DM}}$, in the case of direct (solid), 1-step (dashed), and 2-step (dotted) decays into $e^+e^-$, $\mu^+\mu^-$, and $\tau^+\tau^-$. The best fit values of $m_{\text{DM}}$ and $\tau_{\text{DM}}$ are indicated by the crosses, and are displayed inset in units of TeV and $10^{26}$ sec, respectively. Direct decays into $e^+e^-$ does not give a good fit. The case of $\pi^+\pi^-\pi^0$ is similar to that of $\tau^+\tau^-$. 

marginalize over $r$ in the range $0.9 < r < 1.05$. The result for best fit regions of dark matter mass, $m_{\text{DM}}$, and lifetime, $\tau_{\text{DM}}$ is shown in Figure 4 for 1-step and 2-step cascades to electrons, muons and taus. For comparison, we also show the fits for direct decays. With 26 degrees of freedom (7 PAMELA + 26 FERMI – 7 fitting parameters), we plot 68% CL contours, corresponding to $\chi^2 = 28.8$. The best-fit values of $m_{\text{DM}}$ and $\tau_{\text{DM}}$ are indicated in each case.

We find that good fits are obtained in the region $m_{\text{DM}} \approx O(1 – 100$ TeV) and $\tau_{\text{DM}} \approx O(10^{25} – 10^{27}$ sec), depending on the decay chain. In Figs. 5 and 6 we show the predicted $e^\pm$ fluxes compared to the PAMELA and FERMI data, using the best fit parameters obtained from the $\chi^2$ analysis. The H.E.S.S. high energy data [22] is also overlaid, with the energy rescaled in each case to best match the predicted flux, within the $\pm15\%$ range of overall systematic uncertainty in H.E.S.S. energy. The agreement between the predictions and data is remarkable. We find some tension with the H.E.S.S. data in the case of 1-step to $e^+e^-$, but in all other cases the predicted curves are consistent with the H.E.S.S. data despite the fact that we did not use it in our fit procedure. Note that the background fluxes are very uncertain above $\approx 1$ TeV, so that the precise comparison with the H.E.S.S. data may be misleading; for example, the background spectrum adopted here seems a bit too hard at the highest energies to be consistent with the H.E.S.S. data.

To summarize, we find that the electron/positron fluxes observed by PAMELA and FERMI are very well explained by dark matter decay in our scenario. The mass of the axion-like state can take almost any value in the wide range considered in Section 2.3 except for a small window.
Figure 5: The predicted $e^\pm$ fluxes compared to the PAMELA and FERMI data for 1-step cascade decays into $e^+e^-$, $\mu^+\mu^-$, and $\tau^+\tau^-$. In each case, the mass and lifetime of dark matter are chosen at the best fit point indicated in Figure 4 with the background (dotted) and FERMI energy-normalization marginalized as described in the text. We overlay the H.E.S.S. data with energy rescaled in the range $\pm 15\%$ to best match the theory. Note that due to considerable uncertainty in the background fluxes at H.E.S.S. energies, direct comparison of predicted fluxes with the H.E.S.S. data may be misleading.
Figure 6: The same as Figure 5 for 2-step cascade decays into $e^+ e^-$, $\mu^+ \mu^-$, and $\tau^+ \tau^-$. 
at $\simeq 1.9 - 3.6$ GeV where $a$ decays hadronically. More specifically, we find that the regions

\begin{align}
(i) \quad & a \to e^+e^- \quad (2m_e < m_a < 2m_\mu) & f_a & \approx O(100 \text{ TeV}) , \\
(ii) \quad & a \to \mu^+\mu^- \quad (2m_\mu < m_a \lesssim m_K - m_\pi) & f_a & \approx O(10 - 100 \text{ TeV}) , \\
(ii) \quad & a \to \mu^+\mu^- \quad (m_K - m_\pi < m_a \lesssim 800 \text{ MeV}) & f_a & \approx O(1 - 100 \text{ TeV}) , \\
(iii) \quad & a \to \pi^+\pi^-\pi^0 \quad (800 \text{ MeV} \lesssim m_a < 2m_p) & f_a & \approx O(1 - 100 \text{ TeV}) , \\
(iv) \quad & a \to \tau^+\tau^- \quad (2m_\tau < m_a < 2m_b) & f_a & \approx O(1 - 100 \text{ TeV}) ,
\end{align}

(can all explain the cosmic ray $e^\pm$ data, without conflicting the bounds discussed in Section 2.5).

4.2 Diffuse gamma ray signals at FERMI

An immediate consequence of the framework described here is that the decay of dark matter will provide a source of $\gamma$ rays throughout the dark matter halo, and extending up to energies around a few TeV. If the axion-like state decays into $\tau^+\tau^-$ or $\pi^+\pi^-\pi^0$, photons are produced directly by the decay of $\pi^0$s, while for $e^+e^-$ and $\mu^+\mu^-$ modes there is a much smaller but significant source of $\gamma$ rays from final state radiation (FSR) and inverse Compton scattering (ICS). Although the greatest fluxes would originate from the galactic center, where number densities are highest, the best direction to look for them is away from the galactic plane, where the background is much smaller and the signal still large. The FERMI experiment will measure $\gamma$-ray fluxes over the entire sky at energies up to several hundred GeV, which has the potential to resolve spectral features caused by dark matter decays \cite{64, 65}.

Figure 7: The high energy diffuse $\gamma$-ray spectrum away from the galactic plane, averaged over galactic latitudes above $10^\circ$ assuming an NFW profile. Shown are the best fit parameters for 1-step to $\tau^+\tau^-$ (dashed), 1-step to $\pi^+\pi^-\pi^0$ (dot-dashed), and 2-step to $\mu^+\mu^-$ (dotted) dark matter decay modes. The $\gamma$-rays are due to $\pi^0$ decay in the case of 1-step to $\tau^+$ and $\tau^-$, and to FSR in the case of 2-step to $\mu^+$ and $\mu^-$. The $\tau^+\tau^-$ and $\pi^+\pi^-\pi^0$ modes produce a bump in the flux clearly distinguishable from the background.
In Figure 7 we plot the contributions to the diffuse $\gamma$-ray fluxes for three illustrative decay modes: 1-step to $\tau^+\tau^-$, 1-step to $\pi^+\pi^-\pi^0$, and 2-step to $\mu^+\mu^-$. The first two cases are representative of any cascade to taus or pions, while the $\mu^+\mu^-$ curve illustrates the much lower flux for a decay mode with only FSR photons. Assuming an NFW profile, we average over all galactic latitudes greater than $10^\circ$ from the galactic plane, representative of a diffuse $\gamma$-ray measurement by FERMI. The astrophysical background is modeled by a power law flux, $\propto E^{-2.7}$, taken from [66]. For the $\mu^+\mu^-$ mode the flux shown is solely due to FSR, and we assume an $a$ mass of 600 MeV. We have not shown the contribution from ICS, which, like FSR, is subdominant to directly produce photons for $\tau^+\tau^-$ and $\pi^+\pi^-\pi^0$ modes. At its peak, the ICS flux-component is expected to be comparable to the background, with a spectral shape somewhat different to that from FSR [55, 65], and could be significant for $e^+e^-$ and $\mu^+\mu^-$ modes.

We find that signals from dark matter in our scenario can be seen in the diffuse $\gamma$-ray data. While the background is very uncertain, it is expected to be smooth compared to the strong, peaked fluxes seen for the $\tau^+\tau^-$ and $\pi^+\pi^-\pi^0$ decay modes, which should result in a clearly visible bump in the spectrum. The weaker signals from $e^+e^-$ and $\mu^+\mu^-$ modes may also be seen. Since astrophysical sources are not expected to produce $\gamma$-ray fluxes with such prominent spectral features, measurements of diffuse $\gamma$-rays may serve to distinguish dark matter as the source of the PAMELA and FERMI excesses. Additionally, the shape and the strength of the $\gamma$-ray spectrum would convey information about the mass and decay channels of the dark matter. While the absence of an excess in upcoming FERMI data would not exclude our scenario, especially for the $e^+e^-$ and $\mu^+\mu^-$ cases, a positive result would be a striking signature.

5 Collider Signatures

If the PAMELA and FERMI data is indeed indicative of dark matter with an $O(10 \text{ TeV})$ mass, then direct production of dark matter is not possible at current collider experiments. However, the light axion-like states, which play a crucial role in determining the cosmic ray spectra, are kinematically accessible. We briefly describe some of the collider signatures for the axion-like state $a$.

The discovery potential for the axion-like state depends strongly on its mass and decay constant. Since we are considering decay constants in the range $f_a \approx O(1 – 100 \text{ TeV})$, the coupling of $a$ to the standard model is small, so discovery relies on the axion-like state having a clean decay mode. For $2m_e < m_a < 2m_\mu$, the decay length is generically greater than a kilometer, so any axion-like state produced in a collision will decay outside of the detector. In mass range $2m_\mu < m_a \lesssim 800 \text{ MeV}$, $a$ decays into $\mu^+\mu^-$ with a possible displaced vertex, offering a promising discovery channel. For $2m_\tau < m_a < 2m_b$, $a$ decays promptly into taus, but neither leptonic tau
nor hadronic taus are particularly clean channels. Similarly, there are large hadronic backgrounds to $a \to \pi^+\pi^-\pi^0$. Interestingly, while the $\tau^+\tau^-$ and $\pi^+\pi^-\pi^0$ channels are challenging in the collider context, they are precisely the ones that give the most striking diffuse $\gamma$-ray signal as seen in Section 4.2. In contrast, the $\mu^+\mu^-$ final state yields less dramatic diffuse $\gamma$-ray flux, but relatively clean collider signatures.

Direct production of $a$’s at the LHC was considered in Ref. [24], where the $a$ is produced in association with a hard ($p_T > 100$ GeV) jet via the gluon-gluon-$a$ coupling. Since the $a$ is boosted from the production, the resulting muon tracks from $a$ decay have a mrad opening angle. If the decay constant $f_a \gtrsim (3 - 10)$ TeV, then the axion-like state lives sufficiently long that the decay $a \to \mu^+\mu^-$ happens with an $O$(cm) displaced vertex. As long as the decay constant $f_a < \sim (15 - 30)$ TeV, then the production rate is sufficiently large to compete with the $B^0 \to D^\pm\mu^\pm\nu \to \mu^+\mu^-\nu\nu$ background. Therefore, direct production of $a$ is a promising possibility for $f_a \sim 10$ TeV.

An alternative production mechanism for $a$’s is via the Higgs boson. For the $R$ and PQ symmetries considered in Section 2.3, the axion-like state has couplings to the Higgs fields, allowing the Higgs boson to decay into two axion-like states. For simplicity, we will focus on the decoupling limit as in Eq. (22). The Higgs decay width into two $a$’s is then given by

$$\Gamma(h \to aa) = \frac{c_1^2 v^2 m_h^2}{64\pi f_a^4},$$

where $c_1 = (1/4) \sin^2 2\beta$, and $m_h$ is the Higgs boson mass. For a light Higgs boson, the dominant decay mode is into $b\bar{b}$ with (taking $m_b/m_h \to 0$)

$$\Gamma(h \to b\bar{b}) = \frac{3}{16\pi} m_h \left(\frac{m_h}{v}\right)^2.$$  

Therefore, the branching ratio into two axion-like states is

$$\text{Br}(h \to aa) \approx \frac{c_1^2 m_h^2 v^4}{12 m_b^2 f_a^4} \approx 1.3 \times 10^{-5} \left(\frac{m_h}{120 \text{ GeV}}\right)^2 \left(\frac{10 \text{ TeV}}{f_a \tan \beta}\right)^4,$$

where we have used the large $\tan \beta$ approximation, $\sin 2\beta \approx 2/\tan \beta$, in the last equation. We see that the branching fraction for $h \to aa$ depends strongly on the decay constant as $f_a^{-4}$, and a sizable branching fraction can be obtained as $f_a$ approaches 1 TeV. Since the decay constant is naturally expected to be somewhat ($\approx 4\pi$) smaller than the dynamical scale $\Lambda$ as we saw in Eq. (14), this makes observation of the $h \to aa \to 4\mu$ signal at the LHC an interesting possibility [18]. The recent analysis of Ref. [67] focused on the $h \to aa \to 2\mu 2\tau$ channel in the case of $m_a \sim 5$ GeV. They find that generic event selection has an efficiency around 10%. Assuming a similar efficiency for $h \to aa \to 4\mu$ events, and taking the Higgs production at the
LHC of $\sim 50$ pb, at least 10 events could be seen with 300 fb$^{-1}$ of data for $f_a \tan \beta \lesssim 10$ TeV. In this way, direct $a$ production (for larger values of $f_a$) and $a$ production through the Higgs boson (for smaller values of $f_a$) are complementary. For larger $m_a$, the $h \rightarrow aa \rightarrow 6\pi$ and $h \rightarrow aa \rightarrow 4\tau$ channels may be visible if the background can be controlled. Note that the recent DØ analysis in Ref. [68] already gives the constraint $f_a \tan \beta \gtrsim 2$ TeV for $a \rightarrow \mu^+\mu^-$. There has also been recent interest in looking for light bosons in low-energy high-luminosity lepton colliders [69, 70] as well as in fixed-target experiments [70]. From Eq. (19), the coupling of the axion-like state to electrons is proportional to $m_e/f_a$, which is smaller than $10^{-6}$ for $f_a > 1$ TeV, making the process $e^+e^- \rightarrow \gamma a$ beyond the reach of current lepton colliders. The feasibility of a fixed-target experiment to discover $a$ depends on the $a$ lifetime. If $a$ decays promptly, then one must contend with a huge standard model background of prompt charged particle production, though Ref. [70] suggests that a coupling as small as $10^{-6}$ might be testable if $a$ decays into $\mu^+\mu^-$. If $a$ decays with a displaced vertex, then it could be discovered in traditional electron or proton beam dump experiments. Given the bounds from Figure 2 and the lifetimes in Eqs. (27, 28, 29), the possible values of $c\tau$ spans a huge range from tens of kilometers to less than a nanometer. If $a \rightarrow \mu^+\mu^-$, then $c\tau$ is plausibly in the millimeter to centimeter range, and could likely be discovered in an upgraded version of the experiment from Ref. [38].

6 Discussion and Conclusions

The origin of dark matter is one of the greatest mysteries in particle physics and cosmology. From the theoretical point of view, attentions have been focused on the WIMP paradigm: dark matter has mass $m_{DM} \approx O(100$ GeV – 1 TeV) and couplings of weak interaction strength $g \approx O(1)$, leading to an annihilation cross section that gives the correct thermal relic abundance, $\Omega_{DM} \simeq 0.2$. However, since the annihilation cross section depends on the combination $g^2/m_{DM}$, it should be equally convincing to consider the case where $m_{DM}$ is heavier than the weak scale as long as the coupling $g$ is larger. Such a situation arises naturally if dark matter is a composite state of some strong interaction, since the typical coupling $g$ is expected to be larger than order unity and $m_{DM}$ can be $O(10$ TeV).

Suppose that dark matter indeed arises from some strongly interacting sector. Then its stability may be the result of compositeness, and not of some exact symmetry imposed on the theory. This is precisely analogous to the case of the proton in the standard model embedded in some unified theory. If the proton (and pions) were elementary, it would immediately decay into $e^+$ and $\pi^0$ through a Lagrangian term $\mathcal{L} \sim pe\pi^0$. Since it is composite, however, the leading operator causing proton decay is already dimension six, $\mathcal{L} \sim qql/M_*^2$, and the resulting lifetime is of order $10^{36}$ years for $M_*$ of order the unification scale. In the case of composite dark matter, dimension
six operators yield a lifetime of order $10^{25}$ sec for $m_{\text{DM}} \approx 10$ TeV and $M_* \approx 10^{17}$ GeV. Decay of galactic dark matter could then have currently observable consequences.

The story just described implies that there is new strongly coupled physics beyond the weak scale, at $\approx O(10 - 100$ TeV). Interestingly, we already know an attractive framework where such a situation occurs—weak scale supersymmetry with low energy dynamical supersymmetry breaking. Since the superparticle masses in this framework arise at loop level, the scale of the strong sector is naturally larger than the weak scale by a one-loop factor $\approx 16\pi^2$. This picture is very much consistent with what is implied by the LEP precision electroweak data, namely that physics at the weak scale itself is weakly coupled. Yet such a setup can still explain the large hierarchy between the weak and Planck scales in, arguably, the simplest manner via dimensional transmutation.

It is interesting that this picture of dark matter arises precisely in the scenario where conventional LSP dark matter is lost—the LSP is now the very light gravitino. In fact, a gravitino with mass $m_{3/2} \lesssim O(10$ eV), implied by $\Lambda \approx O(10 - 100$ TeV), avoids various cosmological difficulties faced by other supersymmetry breaking scenarios. This allows us to consider the “standard” cosmological paradigm, based on inflation at a very early epoch with subsequent baryogenesis at high energies, consistently with supersymmetry. The standard virtues of weak scale supersymmetry, such as the stability of the weak scale and gauge coupling unification, are all preserved.

While composite dark matter in low-scale supersymmetry breaking already offers a consistent picture, this may not be the end of the story. Since dark matter is a part of a strongly interacting sector, it is possible that it feels additional dynamical effects. In particular, it is quite plausible that the sector spontaneously breaks an accidental global symmetry, leading to a light pseudo Nambu-Goldstone boson—again as in QCD. This raises the possibility that dark matter decays mainly into these light states (possibly through some other state), which then decay into standard model particles. In supersymmetric theories there is a natural candidate for such symmetry, an $R$ symmetry, whose existence is suggested by a certain genericity argument associated with supersymmetry breaking. The mass of these light states can easily be in the MeV to 10 GeV range if explicit breaking arises from dimension five operators. Except in the mass range $m_a \approx 1.9 - 3.6$ GeV, the decay products of the light states are naturally “leptonic” (specifically, $e^+e^-$, $\mu^+\mu^-$, $\pi^+\pi^-\pi^0$, or $\tau^+\tau^-$), with little nucleonic activity.

An illustrative model of this class was given in Section 3, and the required symmetry structure for the supersymmetry breaking sector was summarized in Table 1. Remarkably, the properties we have just obtained (with a little adjustment of parameters), are precisely what is needed to explain the recent cosmic ray data through dark matter physics:

- Dark matter mass is of $O(10$ TeV),
- Dark matter lifetime is of $O(10^{26}$ sec),
Dark matter decays through (long) cascades,

- Dark matter decay final states are leptonic.

As we saw in Section 4, various mysterious features in the data—an unexpected rise of the positron fraction between \( \approx 10 \) and 100 GeV in the PAMELA data, nonobservation of any anomaly in the PAMELA antiproton data, and a broad excess of the combined \( e^\pm \) flux in the FERMI data—are all beautifully explained by the properties of dark matter discussed in this paper. The resulting \( e^\pm \) spectra are also consistent with the recent H.E.S.S. result. The success is quite remarkable, especially in view of the fact that the data is difficult to explain in terms of conventional WIMP annihilation—the mass scale suggested does not seem natural, the observed rate requires a large boost factor, and typical WIMP annihilation produces more antiprotons than indicated by PAMELA.

While dark matter with mass of order 10 TeV will not be produced at the LHC, the present scenario still has potential collider signatures. Since the light states generically have interactions with standard model gauge bosons, they may be produced at the LHC, and the leptonic decay of the light state could lead to visible signatures, especially if the dominant decay is into muons. The Higgs boson may also decay into the light states, producing a clean four lepton final state. Future astrophysical observations could also probe this scenario. For example, measurements of diffuse \( \gamma \)-ray could discriminate \( e/\mu \) final states from \( \tau/\pi \) final states. Moreover, the present framework may explain the discrepancy between the measured \( ^7\text{Li} \) abundance and the standard big-bang nucleosynthesis prediction by dimension five decay of some of the states in the supersymmetry breaking sector, e.g. (composite) messenger fields. Analysis of all these experimental data could provide important information about the structure of the supersymmetry breaking sector.

We finally mention that while in this paper we focused on the case where dark matter arises from a strongly interacting supersymmetry breaking sector and decays through light states, some of our results apply in wider contexts. For example, quasi-stable dark matter in the supersymmetry breaking sector may decay directly into the SSM sector particles through dimension six operators. For example, quasi-stable mesons can decay into SSM (s)leptons \( L_{\text{SSM}} \) through the Kähler potential operators \( K \sim Q_i^\dagger Q^j L_{\text{SSM}} L_{\text{SSM}} / M_x^2 + \bar{Q}_i^{\dagger} \bar{Q}_j L_{\text{SSM}}^\dagger L_{\text{SSM}} / M_x^2 \). In these cases, the final states of the decay are determined by a combination of gravitational scale physics and TeV-scale superparticle spectra, not through cascades associated with light states. However, all other points regarding compositeness, long lifetime, relatively large mass, and thermal abundance still persist. The dynamics of dark matter discussed here can also be applied in non-supersymmetric theories—all we need is some strong dynamics at \( \approx O(10-100 \text{ TeV}) \) satisfying the properties of Table 1. Examples of such theories may include ones in which the Higgs fields arise as pseudo Nambu-Goldstone bosons of some strong dynamics.

As the LHC starts running this year, we anticipate great discoveries. A possible connection
between weak scale physics and dark matter is one of the major themes to be explored at the high energy frontier. The standard expectation is to study the properties of a dark matter particle by producing it at the LHC. Although the scenario presented in this paper does not allow this, important physics associated with the dark sector may still be probed. With many new particle physics and astrophysics observations on the horizon, the next decade will certainly be exciting for potentially understanding the origin and properties of dark matter. While nature may not show us the “standard” WIMP dark matter story, it may still reveal a beautiful connection between dark matter and the weak scale through hidden sector dynamics.

Acknowledgments

This work was supported in part by the Director, Office of Science, Office of High Energy and Nuclear Physics, of the US Department of Energy under Contract DE-AC02-05CH11231, and in part by the National Science Foundation under grant PHY-0457315. The work of Y.N. was supported in part by the National Science Foundation under grant PHY-0555661 and the Alfred P. Sloan Foundation. J.T. is supported by the Miller Institute for Basic Research in Science.

References

[1] S. Weinberg, Phys. Rev. Lett. 48, 1303 (1982); M. Y. Khlopov and A. D. Linde, Phys. Lett. B 138 (1984) 265; M. Kawasaki, K. Kohri, T. Moroi and A. Yotsuyanagi, Phys. Rev. D 78, 065011 (2008) [arXiv:0804.3745 [hep-ph]].

[2] G. D. Coughlan, W. Fischler, E. W. Kolb, S. Raby and G. G. Ross, Phys. Lett. B 131, 59 (1983); T. Banks, D. B. Kaplan and A. E. Nelson, Phys. Rev. D 49, 779 (1994) [arXiv:hep-ph/9308292].

[3] T. Moroi, H. Murayama and M. Yamaguchi, Phys. Lett. B 303, 289 (1993).

[4] M. Kawasaki, F. Takahashi and T. T. Yanagida, Phys. Lett. B 638, 8 (2006) [arXiv:hep-ph/0603265]; Phys. Rev. D 74, 043519 (2006) [arXiv:hep-ph/0605297].

[5] H. Pagels and J. R. Primack, Phys. Rev. Lett. 48, 223 (1982).

[6] M. Viel, J. Lesgourgues, M. G. Haehnelt, S. Matarrese and A. Riotto, Phys. Rev. D 71, 063534 (2005) [arXiv:astro-ph/0501562].

[7] M. Fukugita and T. Yanagida, Phys. Lett. B 174, 45 (1986).

[8] M. Dine and W. Fischler, Phys. Lett. B 110, 227 (1982); Nucl. Phys. B 204, 346 (1982); L. Alvarez-Gaumé, M. Claudson and M. B. Wise, Nucl. Phys. B 207, 96 (1982); S. Dimopoulos and S. Raby, Nucl. Phys. B 219, 479 (1983).
[9] M. Dine, A. E. Nelson and Y. Shirman, Phys. Rev. D 51, 1362 (1995) [arXiv:hep-ph/9408384]; M. Dine, A. E. Nelson, Y. Nir and Y. Shirman, Phys. Rev. D 53, 2658 (1996) [arXiv:hep-ph/9507378].

[10] C. Csáki, A. Falkowski, Y. Nomura and T. Volansky, Phys. Rev. Lett. 102, 111801 (2009) [arXiv:0809.4492 [hep-ph]].

[11] Z. Komargodski and N. Seiberg, JHEP 0903, 072 (2009) [arXiv:0812.3900 [hep-ph]].

[12] H. P. Nilles, M. Srednicki and D. Wyler, Phys. Lett. B 120, 346 (1983); J.-M. Frère, D. R. T. Jones and S. Raby, Nucl. Phys. B 222, 11 (1983); J.-P. Derendinger and C. A. Savoy, Nucl. Phys. B 237, 307 (1984).

[13] S. Dimopoulos, G. F. Giudice and A. Pomarol, Phys. Lett. B 389, 37 (1996) [arXiv:hep-ph/9607225]; Y. Nomura and B. Tweedie, Phys. Rev. D 72, 015006 (2005) [arXiv:hep-ph/0504246]; K. Hamaguchi, S. Shirai and T. T. Yanagida, Phys. Lett. B 654, 110 (2007) [arXiv:0707.2463 [hep-ph]].

[14] A. E. Nelson and N. Seiberg, Nucl. Phys. B 416, 46 (1994) [arXiv:hep-ph/9309299].

[15] J. Bagger, E. Poppitz and L. Randall, Nucl. Phys. B 426, 3 (1994) [arXiv:hep-ph/9405345].

[16] O. Adriani et al. [PAMELA Collaboration], Nature 458, 607 (2009) [arXiv:0810.4995 [astro-ph]].

[17] I. Cholis, L. Goodenough and N. Weiner, arXiv:0802.2922 [astro-ph]; N. Arkani-Hamed, D. P. Finkbeiner, T. R. Slatyer and N. Weiner, Phys. Rev. D 79, 015014 (2009) [arXiv:0810.0713 [hep-ph]].

[18] Y. Nomura and J. Thaler, Phys. Rev. D 79, 075008 (2009) [arXiv:0810.5397 [hep-ph]].

[19] X. Chen, arXiv:0902.0008 [hep-ph].

[20] E. Nardi, F. Sannino and A. Strumia, JCAP 0901, 043 (2009) [arXiv:0811.4153 [hep-ph]]; A. Arvanitaki, S. Dimopoulos, S. Dubovsky, P. W. Graham, R. Harnik and S. Rajendran, arXiv:0812.2075 [hep-ph], and references therein.

[21] A. A. Abdo et al. [The Fermi LAT Collaboration], arXiv:0905.0025 [astro-ph.HE].

[22] F. Aharonian et al. [H.E.S.S. Collaboration], Phys. Rev. Lett. 101, 261104 (2008) [arXiv:0811.3894 [astro-ph]].

[23] F. Aharonian et al. [H.E.S.S. Collaboration], arXiv:0905.0105 [astro-ph.HE].

[24] H.-S. Goh and M. Ibe, JHEP 0903, 049 (2009) [arXiv:0810.5773 [hep-ph]].

[25] P. Meade, N. Seiberg and D. Shih, arXiv:0801.3278 [hep-ph].
[26] S. Dimopoulos, M. Dine, S. Raby and S. D. Thomas, Phys. Rev. Lett. 76, 3494 (1996) [arXiv:hep-ph/9601367].

[27] E. Witten, Nucl. Phys. B 188, 513 (1981).

[28] K. Hamaguchi, E. Nakamura, S. Shirai and T. T. Yanagida, arXiv:0811.0737 [hep-ph].

[29] J. L. Jones, arXiv:0812.2106 [hep-ph].

[30] M. A. Luty, Phys. Rev. D 57, 1531 (1998) [arXiv:hep-ph/9706235]; A. G. Cohen, D. B. Kaplan and A. E. Nelson, Phys. Lett. B 412, 301 (1997) [arXiv:hep-ph/9706275].

[31] M. Dine, W. Fischler and M. Srednicki, Phys. Lett. B 104, 199 (1981); A. P. Zhitnitskii, Sov. J. Nucl. Phys. 31, 260 (1980) [Yad. Fiz. 31, 497 (1980)].

[32] L. J. Hall and T. Watari, Phys. Rev. D 70, 115001 (2004) [arXiv:hep-ph/0405109].

[33] I. Antoniadis and T. N. Truong, Phys. Lett. B 109, 67 (1982).

[34] V. V. Anisimovsky et al. [E949 Collaboration], Phys. Rev. Lett. 93, 031801 (2004) [arXiv:hep-ex/0403036]; S. Adler et al. [E787 Collaboration], Phys. Rev. Lett. 88, 041803 (2002) [arXiv:hep-ex/0111091].

[35] H. K. Park et al. [HyperCP Collaboration], Phys. Rev. Lett. 88, 111801 (2002) [arXiv:hep-ex/0110033].

[36] B. Aubert et al. [BABAR Collaboration], arXiv:0905.4539 [hep-ex].

[37] W. Love et al. [CLEO Collaboration], Phys. Rev. Lett. 101, 151802 (2008) [arXiv:0807.1427 [hep-ex]].

[38] F. Bergsma et al. [CHARM Collaboration], Phys. Lett. B 157, 458 (1985).

[39] M. Altmann, Y. Declais, F. v. Feilitzsch, C. Hagner, E. Kajfasz, and L. Oberauer, Z. Phys. C 68, 221 (1995).

[40] G. G. Raffelt, Phys. Rept. 198, 1 (1990).

[41] J. Engel, D. Seckel and A. C. Hayes, Phys. Rev. Lett. 65, 960 (1990); E. L. Chupp, W. T. Vestrand and C. Reppin, Phys. Rev. Lett. 62, 505 (1989).

[42] For a review, G. F. Giudice and R. Rattazzi, Phys. Rept. 322, 419 (1999) [arXiv:hep-ph/9801271].

[43] For a review, K. Intriligator and N. Seiberg, Nucl. Phys. Proc. Suppl. 45BC, 1 (1996) [arXiv:hep-th/9509066].

[44] K. Griest and D. Seckel, Phys. Rev. D 43, 3191 (1991).

[45] K.-I. Izawa, F. Takahashi, T. T. Yanagida and K. Yonekura, arXiv:0902.3854 [hep-th].
[46] Z. Komargodski and D. Shih, JHEP 0904, 093 (2009) [arXiv:0902.0030 [hep-th]].

[47] K.-I. Izawa, Y. Nomura, K. Tobe and T. Yanagida, Phys. Rev. D 56, 2886 (1997) arXiv:hep-ph/9705228.

[48] A. Arvanitaki, S. Dimopoulos, S. Dubovsky, P. W. Graham, R. Harnik and S. Rajendran, in Ref. [20]; S. Bailly, K. Jedamzik and G. Moultau, arXiv:0812.0788 [hep-ph], and references therein.

[49] O. Adriani et al. [PAMELA Collaboration], Phys. Rev. Lett. 102, 051101 (2009) arXiv:0810.4994 [astro-ph].

[50] J. Chang et al. [ATIC Collaboration], Nature 456, 362 (2008).

[51] S. Torii et al. [PPB-BETS Collaboration], arXiv:0809.0760 [astro-ph].

[52] See, for example, M. Cirelli, M. Kadastik, M. Raidal and A. Strumia, Nucl. Phys. B 813, 1 (2009) arXiv:0809.2409 [hep-ph].

[53] See, for example, S. Profumo, arXiv:0812.4457 [astro-ph].

[54] See, for example, L. Bergström, J. Edsjo and G. Zaharijas, arXiv:0905.0333 [astro-ph.HE]; D. Grasso et al., arXiv:0905.0636 [astro-ph.HE].

[55] P. Meade, M. Papucci, A. Strumia and T. Volansky, arXiv:0905.0480 [hep-ph].

[56] M. Ibe, Y. Nakayama, H. Murayama and T. T. Yanagida, arXiv:0902.2914 [hep-ph].

[57] T. Banks and J.-F. Fortin, arXiv:0901.3578 [hep-ph]; T. Banks, J. D. Mason and D. O’Neil, Phys. Rev. D 72, 043530 (2005) arXiv:hep-ph/0506015.

[58] See, for example, G. Bertone, M. Cirelli, A. Strumia and M. Taoso, JCAP 0903, 009 (2009) arXiv:0811.3744 [astro-ph]; P. Meade, M. Papucci and T. Volansky, arXiv:0901.2925 [hep-ph].

[59] J. Mardon, Y. Nomura, D. Stolarski and J. Thaler, arXiv:0901.2926 [hep-ph].

[60] T. Sjöstrand, P. Edén, C. Friberg, L. Lönnblad, G. Miu, S. Mrenna and E. Norrbin, Comput. Phys. Commun. 135, 238 (2001) arXiv:hep-ph/0010017; T. Sjöstrand, S. Mrenna and P. Skands, Comput. Phys. Commun. 178, 852 (2008) arXiv:0710.3820 [hep-ph]].

[61] J. F. Navarro, C. S. Frenk and S. D. M. White, Astrophys. J. 490, 493 (1997) arXiv:astro-ph/9611107.

[62] T. Delahaye, R. Lineros, F. Donato, N. Fornengo and P. Salati, Phys. Rev. D 77, 063527 (2008) arXiv:0712.2312 [astro-ph]]

[63] I. V. Moskalenko and A. W. Strong, Astrophys. J. 493, 694 (1998) arXiv:astro-ph/970124; E. A. Baltz and J. Edsjo, Phys. Rev. D 59, 023511 (1999) arXiv:astro-ph/9808243.
[64] E. A. Baltz et al. [GLAST-LAT Collaboration], JCAP 0807, 013 (2008) [arXiv:0806.2911 [astro-ph]].

[65] M. Regis and P. Ullio, arXiv:0904.4645 [astro-ph.GA].

[66] L. Bergström, P. Ullio and J. H. Buckley, Astropart. Phys. 9, 137 (1998) [arXiv:astro-ph/9712318].

[67] M. Lisanti and J. G. Wacker, arXiv:0903.1377 [hep-ph].

[68] V. M. Abazov et al. [DØ Collaboration], arXiv:0905.3381 [hep-ex].

[69] B. Batell, M. Pospelov and A. Ritz, arXiv:0903.0363 [hep-ph]; R. Essig, P. Schuster and N. Toro, arXiv:0903.3941 [hep-ph]; P.-f. Yin, J. Liu and S.-h. Zhu, arXiv:0904.4644 [hep-ph].

[70] M. Reece and L.-T. Wang, arXiv:0904.1743 [hep-ph].

[71] D. B. Kaplan and H. Georgi, Phys. Lett. B 136, 183 (1984); N. Arkani-Hamed, A. G. Cohen and H. Georgi, Phys. Lett. B 513, 232 (2001) [arXiv:hep-ph/0105239]; R. Contino, Y. Nomura and A. Pomarol, Nucl. Phys. B 671, 148 (2003) [arXiv:hep-ph/0306259].