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Dark discrete gauge symmetries

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We investigate scenarios in which dark matter is stabilized by an Abelian Z_N discrete gauge symmetry. Models are surveyed according to symmetries and matter content. Multicomponent dark matter arises when N is not prime and Z_N contains one or more subgroups. The dark sector interacts with the visible sector through the renormalizable kinetic mixing and Higgs portal operators, and we highlight the basic phenomenology in these scenarios. In particular, multiple species of dark matter can lead to an unconventional nuclear recoil spectrum in direct detection experiments, while the presence of new light states in the dark sector can dramatically affect the decays of the Higgs at the Tevatron and LHC, thus providing a window into the gauge origin of the stability of dark matter.

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I. INTRODUCTION

There is compelling empirical evidence for the existence of dark matter (DM) in the Universe [1]. A simple and attractive possibility for DM is a new elementary particle that is stable on cosmological time scales. This stability strongly suggests the existence of a new symmetry in the dark sector, and understanding the precise nature of this symmetry is of fundamental importance in our quest for a theory of DM.

While there are many possibilities for symmetries that stabilize DM, by far the most common example invoked in the literature is a discrete Z_2 symmetry, e.g. R parity in the minimal supersymmetric standard model (MSSM). Generic parity symmetries are clearly a useful tool in building models of DM, but they often appear *ad hoc*. Indeed, discrete symmetries are not an organizing principle of the standard model (SM), and in general, quantum field theories suffer no theoretical inconsistency once discrete symmetries are abandoned. It has also been argued that global discrete symmetries will be violated by quantum gravitational effects [2]. It is therefore reasonable to ask if discrete symmetries may find their origin as a consequence of a more fundamental principle that may ultimately be responsible for the stability of DM.

Gauge symmetry is just such a principle that may provide the eventual rationale for discrete symmetries. A so-called discrete gauge symmetry emerges as a remnant of a spontaneously broken gauge symmetry [2]. These symmetries are exactly conserved owing to their gauge origin, and thus the lightest state having a nontrivial charge under the discrete symmetry is absolutely stable. In the past, discrete gauge symmetries have been used in a variety of contexts, but most notably to explain or find an alternative to R parity in the MSSM [3]. Very few studies have been solely motivated by explaining DM stability, but see Ref. [4] for a recent example based on a Z_2 remnant from a non-Abelian SU(2) gauge symmetry. See Ref. [5] for earlier work on dark matter models with non-Abelian discrete gauge symmetries.

In this paper we will investigate models of hidden sector DM based on a spontaneously broken $U(1)_D$ gauge symmetry which preserves a discrete Z_N subgroup. While $U(1)_D$ is the minimal choice for the dark sector gauge group, it also affords the possibility of renormalizable interactions of DM with the SM via the kinetic mixing [6] and Higgs portal [7] operators. New phenomena may occur for large Z_N symmetries. This includes multiple stable DM candidates, new direct detection signatures, and novel decays of the SM Higgs boson leading to final states with multiple leptons, jets, and missing energy. These signatures provide a means to probe the gauge origin of DM stability.

II. DISCRETE GAUGE SYMMETRY

Let us begin by reviewing the prototype setup for a gauged Z_N discrete symmetry [2]. Consider a $U(1)_D$ gauge theory containing a dark Higgs field ϕ_N with integer charge N and scalar matter field χ_1 with charge 1, where the notation is that the subscript indicates the charge of the field under $U(1)_D$. In addition to the usual self-conjugate terms, the gauge symmetry allows the following term in the Lagrangian:

$$\Delta \mathcal{L} \propto \phi_N^\dagger \chi_1^N + \text{H.c.} \quad (1)$$

We assume the scalar potential is constructed so that the dark Higgs gets a vacuum expectation value (vev). This spontaneously breaks the gauge symmetry, but the Lagrangian preserves a discrete Z_N subgroup of $U(1)_D$ under which the matter fields transform as

$$\chi_1 \to e^{2\pi i k/N}\chi_1, \qquad k = 0, 1, 2, \ldots N-1. \quad (2)$$

In general, as long as there are no additional sources of gauge symmetry breaking and all gauge anomalies are canceled with appropriate matter content, the remnant discrete symmetry will be exactly conserved. The case $N = 2$ corresponds to the usual parity so often encountered in models of DM, but in this construction there is nothing particularly special about the choice of Z_2, as it is simply a consequence of the relative charges of the Higgs and matter fields in the dark sector which are *a priori* undetermined. Note that depending on the field content and associated gauge and Lorentz charges, there may be no renormalizable terms analogous to Eq. (1) present. In this case there will be accidental global symmetries which may also stabilize DM, at least at the renormalizable level. Of course these accidental symmetries may be violated at the nonrenormalizable level, in contrast to the discrete gauge symmetry.

III. MODELS

We now give a survey of DM models based on an Abelian gauged Z_N symmetry, generalizing Eq. (1) to include fermions, multiple matter fields, and chiral matter. Models are characterized by the field content and corresponding charges. The full Lagrangian contains kinetic terms for the $U(1)_D$ gauge boson V_μ and the dark Higgs boson ϕ_N, as well as an appropriate scalar potential such that ϕ_N acquires a vev $\langle\phi_N\rangle = v'/\sqrt{2}$. This vev breaks the gauge symmetry, leaving a massive vector and a physical scalar dark Higgs fluctuation h'. In addition, there are gauge kinetic mixing and Higgs portal terms that allow interactions with the SM; we will discuss these in detail below. Finally, there will be model-dependent terms related to the additional matter fields which comprise the DM sector. We begin by examining models with minimal matter content, i.e. one dark Higgs and one matter field, leading to three possibilities at the renormalizable level: (1) a Z_2 scalar, (2) a Z_2 Dirac fermion, and (3) a Z_3 scalar. We then turn to multifield and chiral models.

A. Z_2 scalar

The Z_2 scalar DM model contains a charge 2 dark Higgs field ϕ_2 and a charge 1 scalar matter field χ_1. The gauge symmetry permits the term (1) to appear in the Lagrangian:

$$\Delta\mathcal{L} = -\frac{\lambda}{\sqrt{2}}\phi_2^\dagger\chi_1\chi_1 + \text{H.c.}, \tag{3}$$

where a possible phase in the coupling λ may be removed by rotating χ_1. A novel aspect of Z_2 models is that Eq. (3) induces a mass splitting between the real and imaginary components of the matter field, $\chi_1 = (S + iP)/\sqrt{2}$, after $U(1)_D$ is spontaneously broken. Thus, the physical masses of the S and P scalars are $m_{S,P}^2 = m_0^2 \mp \Delta m^2$, where m_0^2 includes universal contributions from a bare mass operator $\chi_1^\dagger\chi_1$ and Higgs portal operators $(\phi_2^\dagger\phi_2)(\chi_1^\dagger\chi_1)$, and $(H^\dagger H)(\chi_1^\dagger\chi_1)$ following gauge symmetry breaking, while $\Delta m^2 \equiv \lambda v'$ comes from Eq. (3). According to Eq. (2) the components of the multiplet are Z_2 odd, $S \to -S$, $P \to -P$, so that for negative λ, S will be the stable DM candidate.

The splitting of the individual components S, P leads to an off-diagonal interaction coming from the gauge kinetic term of the matter field:

$$g_D V_\mu P \overleftrightarrow{\partial}_\mu S, \tag{4}$$

where g_D is the gauge coupling of $U(1)_D$. One well-known consequence of this operator is that scattering with nuclei mediated by vectors V_μ becomes inelastic [8]. Another important effect is that this interaction will inevitably allow the heavier Z_2 partner P to decay to the DM plus light SM fermions which will have important implications for collider physics, as we will discuss in detail later.

B. Z_2 fermion

Consider next the minimal theory of fermionic DM stabilized by a discrete Z_2 gauge symmetry. The model contains a charge 2 dark Higgs ϕ_2, and a vectorlike pair of Weyl fermions ψ_1 and ξ_{-1} with charges $+1$ and -1, respectively. Besides the kinetic terms and Dirac mass term, $\mathcal{L} \supset -m_D(\psi_1\xi_{-1} + \text{H.c.})$, the gauge symmetry permits the following terms analogous to Eq. (1):

$$\Delta\mathcal{L} = -\frac{\lambda_L}{\sqrt{2}}\phi_2^\dagger\psi_1\psi_1 - \frac{\lambda_R}{\sqrt{2}}\phi_2\xi_{-1}\xi_{-1} + \text{H.c.} \tag{5}$$

Like the Z_2 scalar above, there is a splitting between the physical mass eigenstates , which can be written in terms of Majorana fermions $\Psi_\pm$, with masses $m_\pm \simeq m_D \pm (m_L + m_R)/2$. Here we have defined $m_{L,R} \equiv \lambda_{L,R}v'$ and made the assumption $m_D \gg m_{L,R}$ for simplicity. There is discrete Z_2 gauge symmetry that remains, under which both $\Psi_\pm$ are odd, ensuring the stability of the lightest particle. In the limit where $\lambda_L = \lambda_R$, the gauge interaction becomes

$$ig_D V_\mu \bar{\Psi}_+\gamma^\mu\Psi_-. \tag{6}$$

We see that, at least structurally, the Z_2 fermion has many similar features to the Z_2 scalar discussed above, and, in particular, the gauge interaction is off diagonal. We note that both the Z_2 scalar and Z_2 fermion models were considered in the past as models of inelastic DM [8,9], and a similar model with two Dirac fermions was examined in Ref. [10].

C. Z_3 scalar

The final model with minimal field content is the Z_3 scalar DM model. This model contains a charge 3 Higgs field ϕ_3 and a charge 1 scalar matter field χ_1. The gauge symmetries allow the additional term (1)

$$\Delta\mathcal{L} = -\frac{\sqrt{2}}{3!}\lambda\phi_3^\dagger\chi_1\chi_1\chi_1 + \text{H.c.} \tag{7}$$

Unlike the Z_2 models considered above, there is no mass splitting between the real and imaginary components, which clearly must be the case since Z_3 charged fields are necessarily complex. After symmetry breaking, the interactions from (7) become

$$-\frac{1}{3!}\lambda v'\chi_1\chi_1\chi_1 - \frac{1}{3!}\lambda h'\chi_1\chi_1\chi_1 + \text{H.c.} \quad (8)$$

The interaction with the dark Higgs h' leads to the possibility of "semiannihilation," $\chi_1\chi_1 \to \chi_1^* h'$, and may have interesting consequences for the thermal history of χ_1, as emphasized recently in [11] (see also [12]). Note that the interactions in (8), being cubic in χ_1, do not allow for a tree-level DM-nucleus scattering. Direct detection is of course still possible through other portal interactions. For other models of Z_3 DM, see Ref. [13].

D. Multifield Z_N models

With a single matter field, as considered in the models above, the renormalizable possibilities for discrete gauge symmetries are limited to Z_2 and Z_3. Once we allow for more than one matter field in the dark sector, renormalizable Z_N models are possible. Here we will consider several interesting features of these models, illustrating them with examples.

Generically, if N is a prime number then, barring accidental global symmetries, there will be only one DM candidate: the lightest particle with nontrivial Z_N charge.[1] However, if N is not prime, but rather composite (has divisors besides 1 and N), then the Lagrangian will preserve additional discrete symmetries that are subgroups of Z_N, allowing for the possibility of multiple stable species and thus multiple DM candidates [14]. Note that multicomponent DM has been considered for a variety of reasons in the past; see Ref. [15]. To illustrate, consider the simplest case in which this occurs, namely, a Z_4 symmetry, which has Z_2 as a subgroup. A multicomponent DM model may be constructed with the following field content: a charge 4 dark Higgs ϕ_4 and scalar matter fields χ_2 and χ_1 with charges 2 and 1. The Lagrangian contains

$$\Delta\mathcal{L} = -\lambda_1\phi_4^\dagger\chi_2\chi_2 - \lambda_2\phi_4^\dagger\chi_2\chi_1\chi_1 - \lambda_3\chi_2^\dagger\chi_1\chi_1 + \text{H.c.} \quad (9)$$

After symmetry breaking, the first term in Eq. (9) leads to a mass splitting for the real and imaginary components S, P of the field χ_2, analogous to the Z_2 scalar model considered previously in Eq. (3), and we take λ_1 negative so $m_S < m_P$. According to Eq. (2) the theory has a descendant Z_4 discrete symmetry under which both χ_2 and χ_1 are charged, as well as a Z_2 symmetry under which S, P are even and χ_1 is odd. The transformation properties are summarized in Table I. It is clear that χ_1, being the only particle with a nontrivial Z_2 charge, is stable and a DM candidate. However, if S is lighter than twice the mass of χ_1, it will also be stable by virtue of the Z_4 symmetry. Hence, in this case there are potentially two species of DM, though what contribution each makes to the cosmological DM depends on other considerations that we will examine later.

[1] Additional heavier particles with Z_N charge can be stable if all possible decay modes are forbidden by kinematics.

TABLE I. Transformation properties of matter fields in the Z_4 (upper table) and Z_6 (lower table) multifield models. Charges q_i for a field χ_i are defined through their transformation under Z_N: $\chi_i \to \exp(\frac{2\pi i}{N}\cdot q_i)\chi_i$.

Z_4 model	χ_1	χ_2
Z_4	$e^{(2\pi i/4)\cdot 1}$	$e^{(2\pi i/4)\cdot 2}$
Z_2	$e^{(2\pi i/2)\cdot 1}$	1

Z_6 model	χ_1	χ_2	χ_3
Z_6	$e^{(2\pi i/6)\cdot 1}$	$e^{(2\pi i/6)\cdot 2}$	$e^{(2\pi i/6)\cdot 3}$
Z_3	$e^{(2\pi i/3)\cdot 1}$	$e^{(2\pi i/3)\cdot 2}$	1
Z_2	$e^{(2\pi i/2)\cdot 1}$	1	$e^{(2\pi i/2)\cdot 1}$

As a slightly more complicated example, consider a Z_6 model with a charge 6 dark Higgs ϕ_6 and three scalar matter fields χ_1, χ_2, and χ_3 with charges 1, 2, and 3. The Lagrangian contains

$$\Delta\mathcal{L} = -\lambda_1\phi_6^\dagger\chi_3\chi_3 - \lambda_2\phi_6^\dagger\chi_3\chi_2\chi_1 - \lambda_3\phi_6^\dagger\chi_2\chi_2\chi_2 - \lambda_4\chi_3^\dagger\chi_2\chi_1 - \lambda_5\chi_3^\dagger\chi_1\chi_1\chi_1 - \lambda_6\chi_2^\dagger\chi_1\chi_1 + \text{H.c.} \quad (10)$$

In this case, we can decompose $Z_6 \cong Z_3 \times Z_2$, so that there are two distinct cyclic symmetries. The transformation properties of the matter fields under these symmetries are displayed in Table I. We see again that there are potentially two species of DM. For example, if χ_1 is the heaviest matter field, then χ_2 and χ_3 will be stable due to their nontrivial charges under Z_3 and Z_2, respectively. Note that χ_3 will be split into its real and imaginary components due to dark symmetry breaking, and the lightest of these will be stable.

From the first example based on the Z_4 model, we learn that a Z_{p^m} symmetry with p a prime number and m a natural number allows for potentially m stable states, while the example with the Z_6 symmetry shows us that, depending on N, Z_N may be decomposed into a direct product of smaller groups [14]. We may surmise that for a general $N = p_1^{m_1}p_2^{m_2}\cdots p_k^{m_k}$, where p_i is prime ($p_i \neq p_j$ for $i \neq j$) and m_i is natural, Z_N is decomposed into the product group

$$Z_N \cong Z_{p_1^{m_1}} \times Z_{p_2^{m_2}} \times \cdots \times Z_{p_k^{m_k}}. \quad (11)$$

There will be at most $m_1 + m_2 + \cdots + m_k$ stable species for a given Z_N symmetry, though the actual number depends on the field content and the spectrum. A simple recipe to obtain the maximum number of stable states is as follows: for each prime p_i in (11), add fields with $U(1)_D$

charges $Q = \{N/p_i, N/p_i^2, \cdots N/p_i^{m_i}\}$ and order the spectrum $m_{N/p_i} < p_i m_{N/p_i^2} < \ldots < p_i^{m_i - 1} m_{N/p_i^{m_i}}$.

Besides multicomponent DM, another important effect which can occur in multifield Z_N models is mass mixing induced by dark symmetry breaking. To illustrate, consider a model with a charge N dark Higgs, and scalar matter fields χ_Q and χ_{Q+N} with charges Q and $Q+N$. Gauge symmetry allows the following term:

$$\Delta\mathcal{L} = -\frac{\lambda}{\sqrt{2}}\phi_N^\dagger \chi_Q^\dagger \chi_{Q+N} + \text{H.c.}, \tag{12}$$

which induces a mass mixing between the matter fields χ_Q and χ_{Q+N} following dark symmetry breaking. This can be diagonalized with an orthogonal rotation leading to mass eigenstates $\chi_{a,b}$. Consider, for simplicity, the limit in which the bare masses are equal, $m_Q^2 = m_{Q+N}^2 \equiv m^2$, and a small mass splitting induced by the operator in (15) given by $\Delta m^2 = \lambda v \ll m$. The mixing angle in this case is 45°, and the masses are split by a small amount $\delta \sim \Delta m^2/M$. Besides a diagonal gauge interaction between the charged matter fields, this rotation induces an off-diagonal gauge interaction:

$$i g_D V_\mu (\chi_a^\dagger \chi_b^\dagger)\begin{pmatrix} Q+N/2 & N/2 \\ N/2 & Q+N/2 \end{pmatrix} \overleftrightarrow{\partial}_\mu \begin{pmatrix} \chi_a \\ \chi_b \end{pmatrix}. \tag{13}$$

These interactions can lead to important consequences for direct detection experiments, as DM may undergo both elastic and inelastic scattering with nuclei, leading to a distinctive recoil spectrum. Moreover, off-diagonal terms allow for novel decay modes of Z_N partners at colliders. We will explore these issues below.

One final common occurrence in Z_N models is the appearance of nontrivial interactions between matter particles. Cubic and quartic vertices are allowed by gauge invariance so long as the corresponding charges cancel, and these can lead to important effects in DM physics. For example, one may have a Yukawa-type interaction between scalars χ_i and fermion matter fields ψ_i,

$$\Delta\mathcal{L} \supset -\lambda \chi_i \psi_j \psi_k + \text{H.c.}, \tag{14}$$

so long as the charges satisfy $i + j + k = 0$. This interaction, along with similar cubic and quartic scalar matter interactions, can manifest in DM annihilation, affecting the relic abundance and indirect detection prospects, and lead to cascade decays of Z_N partners, given a production mechanism at collider experiments.

E. Chiral models

Lastly, we will consider models in which matter consists of a chiral set of fermions. This means the fermion content is chosen so that (1) there are no gauge anomalies, and (2) a vectorlike mass term is forbidden by gauge symmetry. Indeed, it is a striking fact that each generation of SM matter forms a chiral set, given the gauge group $\mathrm{SU}(3)_c \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$, and so it seems quite plausible that DM could also be chiral.

Given that the dark fermions are charged only under $\mathrm{U}(1)_D$, there are two anomaly conditions to satisfy: the $U(1)_D^3$ anomaly and the $\mathrm{U}(1)_D$-gravitational anomaly. To forbid mass terms, any two charges must not sum to zero. For fermions with charges Q_i, these conditions are simply

$$\sum_i Q_i^3 = 0, \qquad \sum_i Q_i = 0, \qquad Q_i + Q_j \neq 0, \quad \text{for all } i, j. \tag{15}$$

Reference [16] analyzed quite generally the problem of finding fermion charges satisfying the conditions (15) and presented many examples of chiral sets. While there are clearly a vast number of possibilities for models of chiral DM, we will be content here to illustrate a few common features in a simple model. Consider the following chiral set:

$$\psi_{-1}^i, \ \psi_3, \ \psi_4, \ \psi_{-6}^i, \ \psi_7, \tag{16}$$

where $i = 1, 2$. This set contains seven "flavors" of Weyl fermions, and it is straightforward to check that the conditions (15) are satisfied. Besides the fermions in the chiral set, we will need at least one dark Higgs to break the gauge symmetry and to generate masses for the fermions. We choose a charge 7 Higgs field ϕ_7 to accomplish this task. This means that ultimately there will be a discrete Z_7 symmetry that stabilizes the lightest fermion mass eigenstate, but as we will see, there may also be accidental global flavor symmetries which can stabilize multiple candidates.

The Higgs ϕ_7 allows for ψ_3, ψ_4 to marry through the Yukawa operator $\phi_7^\dagger \psi_3 \psi_4 + \text{H.c.}$, to form a Dirac fermion. Furthermore, ψ_{-1}^i, ψ_{-6}^j are also married by the Higgs through the following Yukawa interaction:

$$\Delta\mathcal{L} = -\lambda^{ij}\phi_7 \psi_{-1}^i \psi_{-6}^j + \text{H.c.}, \tag{17}$$

where λ^{ij} is a general matrix. This Yukawa mixing is reminiscent of what occurs in the SM, but we should remember that these fermions are not part of separate "generations." The Yukawa matrix can be diagonalized by separate unitary transformations on ψ_{-1}^i and ψ_{-6}^j, and these unitary matrices are not observable with the field content we have specified so far. In particular, there will be no dark flavor change induced by the gauge interactions, as there is an analogue of the Glashow-Iliopoulos-Maiani mechanism operating in this model.

What about the remaining fermion ψ_7 in the chiral set (16)? This fermion is quite similar to the neutrinos of the SM, as no tree-level mass may be generated with the matter and Higgs content considered so far. How may we generate mass for this fermion? First, let us form the $\mathrm{U}(1)_D$ neutral operator $\phi_7^\dagger \psi_7$ having dimension 5/2 and a spinor Lorentz index. This operator, being analogous to HL of the SM, immediately suggests mechanisms for mass generation. A Majorana mass term may arise from the dimension-five operator $(\phi_7^\dagger \psi_7)^2$, or we may add a singlet fermion ξ_0 and form a Dirac mass $\phi_7^\dagger \psi_7 \xi_0$ + H.c.. As the singlet may also have a bare mass, a seesaw mechanism is potentially operative. If $\mathrm{U}(1)_D$ is broken around the weak scale, and the dimension-five Majorana mass operator or the singlet mass is generated at a higher scale, then typically we are led to predict a much lighter fermionic state in the spectrum. This state is potentially stable and may be difficult to deplete in the early universe, having few annihilation channels which are unsuppressed. There are many ways out of this. For example, the singlet ξ_0 may couple to the SM neutrinos, breaking the accidental flavor symmetry and allowing the state to decay. We may add new dark Higgs fields to make the neutrinolike state heavy, one example being a charge 14 Higgs which would not disturb the subgroup discrete Z_7 symmetry. Alternatively, a very light or massless state in the dark sector could manifest itself as a new component to the cosmological dark radiation and have interesting signatures in its own right.

Finally, let us elaborate on the symmetries of this model. By construction, there is a Z_7 discrete symmetry which will forbid the lightest state with nontrivial charge from decaying. However, there are additional accidental flavor symmetries present which lead to multiple stable states. Let us trace the origin of these symmetries. First, if we turn off all interactions, there is a large U(7) flavor symmetry, which is broken to $\mathrm{U}(2)^2 \times \mathrm{U}(1)^3$ by the gauge interactions. The Yukawa terms further break the flavor symmetry to $\mathrm{U}(1)^4$, leading to potentially four stable mass eigenstates. The U(1) flavor symmetry associated with the neutrinolike state ψ_7 may be broken further to a Z_2 symmetry by the inclusion of a dimension-five Majorana mass operator, or to a diagonal subgroup of the combined $\mathrm{U}(1)^2$ symmetry with the inclusion of the singlet ξ_0. Finally, higher dimension operators will generically further break these accidental symmetries. For example, integrating out a charge 2 scalar field (which does not condense) generates operators like $(\psi_{-1}\psi_{-1})(\psi_{-1}\psi_3)$, $(\psi_4\psi_{-6})(\bar{\psi}_{-1}\bar{\psi}_{-1})$, etc., allowing some of the would-be stable states to decay. However, the discrete Z_7 symmetry is exactly conserved and will ensure the stability of the lightest charged state, thus providing a DM candidate.

IV. PHENOMENOLOGY

We will now describe basic aspects of the phenomenology in this class of models. Detailed scans of the parameter spaces of specific models are beyond the scope of the present work. Rather, we will focus on the novel implications of a larger Z_N discrete symmetry and the impact of additional states in the dark sector such as Z_N matter as well as dark gauge and Higgs bosons. The first question we must address is the following: how can the dark sector communicate with the SM?

A. Portals

Given that we are considering models in which DM is in a hidden sector [DM carries no charge under the SM gauge group, and likewise the SM fields carry no charge under $\mathrm{U}(1)_D$], if interactions are to exist we expect that at high scales there exist heavy states charged under both sectors. Integrating these out, we generate all higher-dimensional operators allowed by the gauge symmetries of the theory. We may also generate "portal" operators [17]: renormalizable, gauge-invariant operators that connect the SM to the dark sector. In particular, a dark sector based on a broken $\mathrm{U}(1)_D$ gauge symmetry affords the opportunity to interact through both the kinetic mixing [6] and Higgs portal [7]:

$$\mathcal{L}_{\mathrm{portal}} = -\frac{\kappa}{2} V_{\mu\nu} B^{\mu\nu} - 2\lambda_\phi \phi^\dagger \phi H^\dagger H - 2\lambda_\chi \chi^\dagger \chi H^\dagger H, \tag{18}$$

where of course the final term is present only for scalar matter fields at the renormalizable level. Note that the combination of vector and Higgs portals was first studied in Ref. [18]. Besides kinetic mixing between vectors, symmetry breaking induces mass mixing between the SM and dark Higgs bosons. We will limit ourselves to the case of small mixing arising from portals, so that it suffices to treat the terms in Eq. (18) as interactions.

Alternatively, one can diagonalize the Lagrangian and work directly with the physical states. The main effect of the kinetic mixing is that SM matter picks up a small charge under $\mathrm{U}(1)_D$ proportional to the kinetic mixing strength κ:

$$\mathcal{L} \supset \kappa V_\mu [-c_w J^\mu_{\mathrm{EM}} + s_w (1 - m_Z^2/m_V^2)^{-1} J^\mu_Z]. \tag{19}$$

The Higgs portal mass mixing may be treated by first expanding around the vacuum, $v \to v + h$ and $v' \to v' + h'$, followed by diagonalizing the system with an orthogonal rotation by the angle $\theta_h \simeq 2\lambda_\phi v v'/m_h^2$, valid in the limit of $m_{h'} \ll m_h$. Diagonalization induces interactions between the SM Higgs and the dark sector, as well as the dark Higgs with the SM,

$$\mathcal{L} \supset \theta_h (h' J_h - h J_{h'}), \tag{20}$$

where $J_{h,h'}$ are the SM and dark currents coupling to the Higgs bosons. Note that for a dark symmetry breaking scale well below the weak scale, there can be an issue with technical naturalness since the SM Higgs vev will tend to raise the dark Higgs vev for large Higgs portal coupling λ_ϕ in Eq. (18). In some cases we will be interested in couplings of order the bottom quark Yukawa, so that e.g. a symmetry breaking scale in the dark sector of $v' \sim 10$ GeV requires only a mild order-one cancellation between different terms in the scalar potential. Scales much lighter than this require smaller values of λ_ϕ for technical naturalness. Similar comments apply to the mass of scalar DM when the Higgs portal coupling λ_χ in Eq. (18) is sizable.

B. Existing constraints

Since the dark sector carries no SM charges, new states can, in principle, be much lighter than the weak scale and still avoid constraints coming from direct production and precision measurements. The constraints that do exist depend on the strength of the portals and the precise decay patterns of dark states. We now give an overview of the possible constraints for various mass scales in the dark sector.

Let us first discuss constraints on the kinetic mixing portal. If the dark gauge boson mass is above the weak scale, there are mild constraints on the kinetic mixing parameter κ coming from electroweak precision tests, $\kappa \lesssim 10^{-1}$–10^{-2} [19]. Vectors with masses $O(1\text{ GeV}) < m_V < O(100\text{ GeV})$ can be constrained further from a variety of e^+e^- collider data, and give a rough bound of $\kappa \lesssim \text{few} \times 10^{-2}$ [20]. For sub-GeV vectors, additional constraints arise from corrections to the anomalous magnetic moment of the electron and muon, extending κ down to $\lesssim \text{few} \times 10^{-3}$ [21]. There are also model-dependent constraints for light vectors, assuming the dominant decay mode of V occurs through the kinetic mixing portal back to SM states, arising from searches at B and other meson factories as well as fixed target experiments [22–25]. In particular, the fixed target experiments are able to exclude a large portion of the parameter space for vector masses below a few hundred MeV, with constraints on κ ranging from 10^{-4}–10^{-8} [24,25]. There has been a great deal of effort invested in the last decade to constrain or search for new light gauge bosons, primarily motivated by possible DM connections [26,27] to astrophysical anomalies [28–30]. There are the additional possibilities of detecting or constraining light gauge bosons decaying to leptons through searches of the existing B-factory data sets [22] and new fixed target experiments [24]. While we have been focused on kinetically mixed gauge bosons [6] in the MeV–100 GeV range, light gauge bosons have a long history and have been considered as long ago as Ref. [31]. For a review considering the constraints from even lighter bosons, see Ref. [32].

Constraints on the Higgs portal are much weaker, and only model-dependent bounds are possible. Very large $O(1)$ mixing angles θ_h can affect electroweak precision observables [33], but for a subweak scale dark Higgs h', a stronger constraint comes from direct production at the CERN LEP experiments. The mixing induces a coupling of the dark Higgs to the Z boson, allowing for a dark Higgs-strahlung process $e^+e^- \to Zh'$. Data from the LEP experiments constrain the parameter $\xi^2 \equiv (g_{\rm HZZ}/g_{\rm HZZ}^{\rm SM})^2 \lesssim 10^{-2}$ [34] for a very light Higgs-like state, which can be translated to a bound on the mixing angle $\theta_h \lesssim 0.1$ which is easily satisfied for $\lambda_\phi \sim y_b$. Note that this bound is conservative and could, in principle, be weakened considerably in specific models depending on the decay modes of h'. For very light states in the GeV range, Higgs portal couplings can be probed by rare heavy-flavor meson decays [35], potentially down to mixing angles of 10^{-2}–10^{-3}, though again, these constraints depend on the decays of the GeV-scale dark states.

C. Relic abundance

Assuming that annihilation processes dominate at freeze-out, the relic abundance of a stable heavy particle is given by $\Omega_{\rm DM} h^2 \simeq 0.1 \times (\langle\sigma v\rangle/\text{pb})^{-1}$, where $\langle\sigma v\rangle$ is the DM annihilation cross section. The predicted relic abundance is to be compared with the observed value from the Wilkinson Microwave Anisotropy Probe, $\Omega_{\rm DM} h^2 = 0.1123 \pm 0.0035$ [36]. We may classify the annihilation channels according to whether the final states belong to the dark sector or the SM:

$$\begin{aligned} \mathbf{I}: &\ \chi + \bar\chi \to VV,\, Vh',\, h'h', \ldots, \\ \text{II}: &\ \chi + \bar\chi \to V^*,\, h'^*,\, h^* \to \psi_{\rm SM} + \bar\psi_{\rm SM}. \end{aligned}$$

Case I corresponds to the "secluded" DM regime [37], in which the relic abundance does not depend on the strength of mediation to the SM. The secluded regime is typically most important when the DM candidate(s) is not the lightest state in the dark sector and direct annihilation channels are kinematically open in the dark sector. For example, if $m_V < m_\chi$ the annihilation cross section is parametrically $\langle\sigma v\rangle_{\chi\bar\chi\to VV} \approx \pi\alpha_D^2/m_\chi^2$, so that the observed Wilkinson Microwave Anisotropy Probe abundance requires $\alpha_D \gtrsim 5 \times 10^{-3}(m_\chi/200\text{ GeV})$ [37].

If no annihilation channels are available directly in the dark sector, then it is still possible to obtain the proper relic abundance provided the portal operators are active and have a sizable coupling. This is case II above, in which DM may annihilate through the kinetic mixing or Higgs portal operators into the light SM final states. As an example, consider the Z_2 scalar model of Eq. (3). If the DM candidate S is the lightest state in the dark sector, it must annihilate into SM states if it is to be thermally depleted. One option is to utilize the "direct" Higgs portal operator in Eq. (18), $(H^\dagger H)(\chi^\dagger\chi) \supset hSS$. In the regime $m_S \ll m_h$ and away from the s-channel resonance at $m_S = m_h/2$, the dominant annihilation mode is into $b\bar b$. The annihilation cross section is $\langle\sigma v\rangle_{SS\to b\bar b} \simeq 3\lambda_\chi^2 m_b^2/\pi m_h^4$, which dictates that $\lambda_\chi \gtrsim 0.2 \times (m_h/120\text{ GeV})^2$. This is a sizable coupling, much larger than the bottom quark Yukawa, and so this model makes the interesting prediction that the SM Higgs decays to DM. In fact, in the limit that all other states in the dark sector are heavy, this model matches on to the minimal Z_2 scalar singlet model of Refs. [7,38,39].

Next, we consider the relic abundance in multicomponent models. Typically we might expect that if two DM particles have distinct annihilation channels governed by different couplings and mass scales, then barring some accidental coincidence in these couplings, the two DM states will end up with dramatically different relic densities and one state will dominantly make up the observed cosmological DM. The situation is somewhat more constrained if the two states have the same annihilation modes and symmetry guarantees the equality of couplings. This occurs, for example, when two heavy states annihilate into light vectors. Given our discussion above regarding this channel, we expect then if two states χ_i, with $i = 1, 2$, have charges Q_i and masses m_i, their relative abundance will be given by $\Omega_1/\Omega_2 = (m_1^2/m_2^2)(Q_2^4/Q_1^4)$. In this case, a similar relic abundance requires a coincidence of masses and charges. However, even if a particular species is subdominant in the cosmic DM, it still may give signatures in direct and indirect detection experiments. We will illustrate this below for multicomponent direct detection.

The final interesting possibility we wish to mention is the process of semiannihilation, examined recently in Ref. [11]. Typical annihilation processes change the number of DM particles by two units. In contrast, semiannihilation reactions change the number by only one unit. The simplest example of this occurs in the Z_3 scalar model in Eq. (7) [11]. If the dark Higgs h' is lighter than the DM χ_1, then the process $\chi\chi \to \chi h'$ will occur with an annihilation cross section $\langle\sigma v\rangle_{\chi\chi\to\chi h'} = 3\lambda^2/128\pi m_\chi^2$, implying that $\lambda \gtrsim 0.1 \times (m_\chi/200\ \text{GeV})$ if other annihilation processes are small. In fact, semiannihilation occurs quite generally in multicomponent DM models, where processes such as $\chi_i\chi_j \to \chi_k h'$, etc., can take place, and we refer the reader to Ref. [11] for more details.

D. Direct detection

We now discuss the possibility of direct detection of DM. Scattering with nuclei is mediated via a t-channel exchange of vectors or Higgs bosons through the portal interactions in (18). Vector exchange leads to the following DM-nucleon scattering cross section [37]:

$$\sigma_n \simeq \frac{16\pi\kappa^2 c_w^2 \alpha\alpha_D Q^2 \mu_n^2}{m_V^4}\frac{Z^2}{A^2}. \tag{21}$$

The most recent limits from the CDMS-II [40] and XENON100 [41] experiments have reached a sensitivity of $\sigma_n \lesssim (3\text{–}4) \times 10^{-44}\ \text{cm}^2$ for DM with masses larger than around 50 GeV [42]. This translates to a bound $\kappa g_D Q \lesssim 10^{-5} \times (m_V/10\ \text{GeV})^4$ for the vector portal.

Exchange of the SM Higgs boson due to the portal interaction also mediates elastic scattering with nuclei. As the simplest example, consider the Z_2 scalar model in Eq. (3). The S-nucleon cross section is (see e.g. [39])

$$\sigma_n \simeq \frac{\lambda_\chi^2 f^2 m_n^4}{\pi m_\chi^2 m_h^4}, \tag{22}$$

where $f \equiv \sum\langle n|m_q\bar{q}q|n\rangle/m_n \simeq 0.35$, with n being a proton or neutron [43]. Constraints from CDMS-II [40] and XENON100 [41] imply $\lambda_\chi \lesssim 0.03 \times (m_\chi/50\ \text{GeV}) \times (m_h/120\ \text{GeV})^2$. Note that for smaller masses in the 10 GeV range, the sensitivity of these experiments falls off rapidly and larger values of the portal couplings can be accommodated.

One possibility to explain the DAMA [44] anomaly is inelastic DM [8]. For a recent study of the status of inelastic DM, see Ref. [45]. As we have seen, mass splittings may occur very easily in particular models with discrete gauge symmetries. For example, the models with a Z_2 discrete symmetry naturally lead to inelastic scattering due to the off-diagonal coupling to the vector V in Eqs. (4) and (6), and in fact, very similar models were originally considered as examples in [8]. More generally, inelastic couplings occur for a multifield Z_N model when there exists matter fields $\chi_{N/2}$ with charge $N/2$, or two matter fields with charges Q and $Q + N$ as in Eq. (12). Small mass splittings of $O(100\ \text{keV})$ can lead to a very long-lived excited state with a significant relic abundance, as pointed out in Ref. [46], and this leads to the possibility of exothermic down-scattering with nuclei. Recently, this has been utilized in Refs. [47,48] as a possible explanation of DAMA.

There is also the possibility of "semielastic" scattering, investigated recently in Ref. [49], which has a very distinctive recoil spectrum. Semielastic scattering can occur due to the exchange of a vector with an off-diagonal coupling and a Higgs with a diagonal coupling, as in [49], or through the exchange of a vector with both diagonal and off-diagonal couplings as in Eq. (13). In the latter case, in order to have the small mass splitting required for an inelastic transition, the bare masses of the two fields must be approximately equal, and we have not provided a symmetry reason for this here, though, in principle, this could be addressed with further model building.

Finally, we turn to the direct detection signatures of multicomponent DM, which has been considered previously in Ref. [50]. Naively, one might expect that detecting two components is only feasible if both DM particles have similar relic densities. As we discussed above, this is not generally true unless both particles have similar masses and annihilation channels. However, the rate depends both on the relic density and on the scattering cross section. A particle with a large annihilation cross section that is efficiently depleted early in the universe might also be expected to have a larger scattering cross section with nuclei. Furthermore, the particle with the dominant relic density can scatter inelastically with nuclei, which leads to a smaller cross section compared to the elastic case. These two facts suggest that it may indeed be possible to have significant rates for both DM particles at direct detection experiments.

As an example, consider the Z_6 model in Eq. (10) with two stable particles, the real component S of χ_3 and χ_2. These states are stable due to the presence of Z_3 and Z_2 symmetries as explained above. Assuming annihilations to

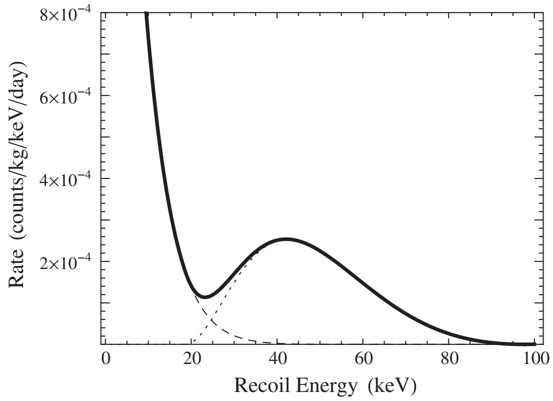

FIG. 1. Recoil spectrum of two-component scattering at XENON100. Here we consider a Z_6 model in Eq. (10) with two DM particles $S \supset \chi_3$ and χ_2. We have assumed masses for the DM particles $m_S = 3$ TeV with a splitting of $\delta \equiv m_P - m_S = 150$ keV leading to inelastic transitions, and $m_{\chi_2} = 30$ GeV. Other parameters entering into the relic abundance and scattering cross section are $\kappa = 8 \times 10^{-4}$, $\alpha_D \sim 0.01$, and $m_V = 10$ GeV. This point is consistent with current direct detection constraints [40,41] and should lead to ∼15 events at XENON100 assuming an effective exposure of 3000 kg-days. The dashed (dotted) lines show the elastic (inelastic) components, while the solid line shows the total rate for the summer (June 2nd). To compute the rate, we have used a standard Maxwellian velocity distribution with $v_{\rm esc} = 500$ km/s.

dark gauge bosons V, whichever particle is heavier will dominantly comprise the cosmological DM, and we assume this to be S. Because the S is split from its partner P, an off-diagonal gauge interaction as in Eq. (4) mediates inelastic scattering with nuclei, and this will be quite suppressed for splittings of $O(100$ keV). On the other hand, χ_2 will undergo elastic transitions, and because of its dilute presence in the galactic halo, we are free to increase κ to much larger values without conflicting with direct detection constraints. If there is a large mass separation between the two DM components, a very distinctive recoil spectrum can be observed at current and next generation experiments, as we illustrate in Fig. 1.

E. Collider signatures

When considering new particles from a dark sector, because of the weak coupling to the SM, it is typically hopeless to directly produce such states at colliders. However, an interesting opportunity arises if a hidden valley scenario is operative [51], in which case we first produce new heavy "connector" states with SM quantum numbers in abundance, which subsequently decay into the dark sector states. While there are many possibilities for connector states around or above the weak scale, we will focus on the minimal case of the SM Higgs boson as our connector. The SM Higgs is in fact an ideal choice for a bridge to the dark sector because of its sizable production cross section, $\sigma^{\rm NNLO}_{gg\to h} \sim 50(10)$ pb for the LHC at $\sqrt{s} = 14(7)$ TeV, and narrow width. For $m_h < 130$ GeV, the Higgs decay width is governed by the small bottom quark Yukawa $y_b \sim 1/60$, so that a light SM Higgs is quite susceptible to new decay modes [52]. We will utilize this property of the Higgs to gain access to states in the dark sector.

From Eq. (18) we see there are two possible operators that will allow for new decays of the SM Higgs. The operator $\phi^\dagger \phi H^\dagger H$ allows for the decay of the SM Higgs to dark Higgs and gauge bosons, with partial widths

$$\Gamma_{h\to h'h'} = \frac{\lambda_\phi^2 v^2}{8\pi m_h}(1 - 4x_{h'}^2)^{1/2}, \tag{23}$$

$$\Gamma_{h\to VV} = \frac{\lambda_\phi^2 v^2}{8\pi m_h}(1 - 4x_V^2)^{1/2}(1 - 4x_V^2 + 12x_V^4), \tag{24}$$

where $x_i \equiv m_i/m_h$. The SM Higgs couples to the dark vectors due to the mass mixing induced by $\mathcal{L}_{\rm portal}$. Notice that in the limit $x_{h',V} \ll 1$, the branching to dark Higgs bosons and dark vectors is equal, $\Gamma_{h\to h'h'} = \Gamma_{h\to VV} \simeq \lambda_\phi^2 v^2/8\pi m_h$, a simple consequence of the Goldstone-equivalence theorem since the Higgs portal (18) leads to identical couplings of the SM Higgs to the radial and Goldstone modes of the dark Higgs. We should compare these partial widths with the partial width to a pair of bottom quarks, $\Gamma_{h\to b\bar{b}} \simeq 3y_b^2 m_h/16\pi$, telling us that for $\lambda_\phi \sim y_b$ the branching into dark Higgs and gauge bosons will be significant. Note that this value of λ_ϕ corresponds to a mixing angle $\theta_h \sim 10^{-3}$–10^{-2}, well below any existing constraints.

If the theory contains scalar matter fields χ, where χ now generically denotes either the stable DM state or an unstable Z_N partner, then the second Higgs portal operator $\chi^\dagger \chi H^\dagger H$ in Eq. (18) allows a decay $h \to \chi^\dagger \chi$, with partial width

$$\Gamma_{h\to \chi^\dagger\chi} = \frac{\lambda_\chi^2 v^2}{4\pi m_h}(1 - 4x_\chi^2)^{1/2}. \tag{25}$$

A similar conclusion applies: as long as λ_χ is comparable to y_b, the SM Higgs will have a sizable branching into DM or its Z_N partners. We show the branching ratios of the SM Higgs in Fig. 2 for a sample model and spectrum.

After the primary decay of the SM Higgs into either $h'h'$, VV, or $\chi^\dagger\chi$, the observed signature depends on the cascade decays of the dark states. This becomes highly dependent on both the field content and the spectrum in the dark sector. Even a small number of additional unstable Z_N matter states in the dark sector can lead to very complicated signatures due to long cascades, particularly when there is mass mixing or nontrivial Z_N interactions as in (14). It is therefore useful to first consider a simplified situation where there is one stable DM χ and a heavier unstable Z_N partner χ^*, in which case we can classify the possible decay patterns as follows:

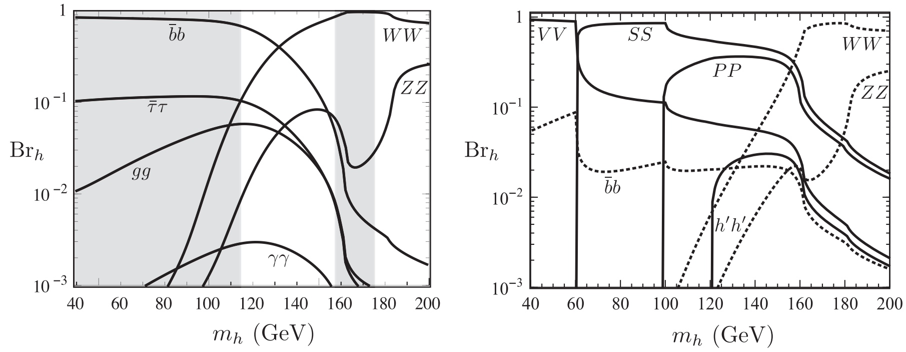

FIG. 2. Higgs branching ratios. The left panel shows the branching ratios of h in the SM along with the direct constraints from the LEP [53] and the combined CDF/D0 constraints [54] in the shaded region. The right panel illustrates the effect of the dark sector. Here we are considering the Z_2 scalar model of Eq. (3), with a spectrum $(m_V, m_S, m_P, m_{h'}) = (20, 30, 50, 60)$ GeV. We have taken portal couplings to be $\lambda_\phi \sim y_b = 1/60$ and $\lambda_\chi = 1/20$. Below the WW threshold, the SM Higgs decays dominantly to hidden sector states.

(i) χ stable ($\not{E}$),
(ii) $\chi^* \to 2\chi, V\chi, h\chi$,
(iii) $V \to 2\chi, 2\chi^*, \chi\chi^*, \ell^+\ell^-, jj$,
(iv) $h' \to 2\chi, 2\chi^*, \chi\chi^*, VV, jj$.

Note that this is schematic, since for the matter states χ, χ^* we have not distinguished between the particle and antiparticle, and furthermore, in particular models the real and imaginary components of matter fields may receive mass splittings. We observe two important points from the decay patterns above: (1) It is quite common to decay into a stable DM state, which will lead to a missing energy signature, and (2) dark states may cascade decay back to the SM, leading to a dilepton or dijet signature. Thus the generic signatures in this class of models are multilepton, multijet final states with missing energy.

While the origin of the missing energy signature is fairly obvious, it is useful to examine in more detail how it is possible for dark sector states to decay back to the SM. This occurs due to the kinetic mixing and Higgs portals in Eq. (18). For the vector, the kinetic mixing operator $V_{\mu\nu}B^{\mu\nu}$ leads to the coupling (19) of V to the SM fermions. The partial width for a vector to decay to a generic SM fermion f is $\Gamma_{V\to f\bar{f}} \simeq (N_c/3)\alpha\kappa^2 m_V$. If there are no direct channels open to dark sector states, the vector will decay to SM fermions. For very light vectors, below the hadronic threshold, the vectors decay into e^+e^- pairs or, if kinematically allowed, $\mu^+\mu^-$ pairs. Above the dipion threshold, the vector can decay to hadrons; the hadronic yield is governed by the well-measured form factor $R = \sigma_{e^+e^-\to\mathrm{hadrons}}/\sigma_{e^+e^-\to\mu^+\mu^-}$ for vectors lighter than several GeV [22,55]. At higher masses, one can use the quark description to calculate the decay to hadrons [56]. In summary, the vector decays fairly democratically, and, in particular, there is always a significant $O(1)$ branching to dilepton final states. Finally, the decays are typically prompt for $\kappa \gtrsim 10^{-5}$, but there may be displaced vertices for smaller kinetic mixing parameters. Similarly, the dark Higgs h' can decay via the Higgs portal to the heaviest kinematically available SM fermion pair. Typically this is the b-quark pair, and so the partial decay is given by $\Gamma_{h'\to b\bar{b}} \simeq 3\theta_h^2 y_b^2 m_h/16\pi$, which is a prompt decay for $\theta_h \gtrsim 10^{-4}$. A leptonic signature from the Higgs portal is not possible unless the dark Higgs is very light, in the hundreds of MeV range, in which case there is some tension with technical naturalness. From now on we will focus on the clean leptonic signature via the kinetic mixing portal. The possibility of the SM Higgs decaying to multilepton final states in this manner was discussed in Ref. [56].

Because of the kinetic mixing suppression in the decay $V \to f\bar{f}$, if kinematically allowed the dark vector will prefer to decay directly into states in the dark sector. However, this does not necessarily imply an absence of SM particles in the final states, since other dark sector states may decay back to the SM through an off-shell V^*. This occurs when mass splittings or mixings lead to off-diagonal gauge interactions, as given e.g. in Eqs. (4), (6), and (13). As an illustration, consider the Z_2 scalar model with the gauge interaction in Eq. (4), and a spectrum such that $m_V \gg m_P > m_S$, so that on-shell vectors decay dominantly via $V \to PS$. The heavy scalar partner will decay via the three-body process $P \to SV^* \to Sf\bar{f}$, and in the limit of large m_V and small fermion mass, the partial width is given by

$$\Gamma_{P\to S\bar{f}f} = \frac{N_c\alpha\alpha_D\kappa^2(g_V^2 + g_A^2)}{48\pi c_w^2}\frac{m_P^5}{m_V^4}f\left(\frac{m_S^2}{m_P^2}\right), \qquad (26)$$

where $f(\hat{x}) = 1 - 8\hat{x} - 12\hat{x}^2\log\hat{x} + 8\hat{x}^3 - \hat{x}^4$. The quantity $(g_V^2 + g_A^2)$ can be extracted from Eq. (19) and depends on the ratio m_V/m_Z; light vectors $m_V/m_Z \ll 1$ couple to electric charge, $(g_V^2 + g_A^2) \simeq c_w^4 Q_{\mathrm{EM}}^2$, while heavy vectors $m_V/m_Z \gg 1$ couple to hypercharge, $(g_V^2 + g_A^2) \simeq (Y_L^2 + Y_R^2)/2$. For a particular fermion final state to be allowed kinematically, the mass splitting must be large enough, $m_P - m_S > 2m_f$. The lifetime depends sensitively on the masses of the vector and Z_2 partner P as well as the kinetic mixing, but for masses in the 10–100 GeV range, the displaced vertices may occur even for large values of the kinetic mixing $\kappa \lesssim 10^{-3}$.

An interesting probe of multicomponent Z_N models would be the presence of two or more distinct stable particles in the cascade decays of the SM Higgs. This may happen from an initial asymmetric decay of the SM Higgs into two distinct states, $h \to \chi_i \chi_j$ with $i \neq j$, which subsequently cascade into distinct stable states plus SM fermions. Even if the initial decay is symmetric, e.g. $h \to VV$, if V has comparable branchings into distinct final states, then there will be events which contain different stable particles. Consider as an example the Z_4 model of Eq. (9) with two stable candidates S and χ_1. Provided the vector is much heavier than the DM candidates, the vector will decay via $V \to SP$ and $V \to \chi_1^\dagger \chi_1$ with partial widths proportional to the square of the charges. Thus, the yield of SP vs $\chi_1^\dagger \chi_1$ will be 4:1. The signature in this case would be $\ell^+ \ell^- \not{E}$, with missing energy coming from two S particles on one side of the event and two χ_1 particles on the other. There are of course many other possibilities and, in general, it would be worthwhile to investigate the extent to which the properties (mass, spin, etc.) of the different stable states can be determined. There has been recent work in this direction in Ref. [57]. Along these lines, an interesting question is whether colliders have something to say regarding the nature of the discrete symmetry at hand; i.e. can one distinguish between, say, Z_2 and Z_3 symmetries? This question has received attention recently in Ref. [58], where the authors study the difference in the possible decays of parent particles obeying a Z_2 vs Z_3 symmetry and how such differences manifest in the invariant mass distributions of the outgoing SM particles. Another avenue of determining the underlying discrete gauge symmetry in these models is to extract the charges of the Higgs boson and matter fields from collider data. Measuring specific decay widths and branching ratios of e.g. the gauge boson into matter states can provide some information in this regard.

If the dark sector is very light compared to m_h, say at a scale of O(GeV) or less, the decay products will be highly boosted, leading to a novel "lepton jet" signature [59]: multiple highly collimated groups of leptons in the final state. In fact, it has recently been suggested that a SM Higgs decaying to lepton jets could have been missed at the LEP experiments [60], allowing for a relaxed bound on the SM Higgs mass. As in the unboosted case, various aspects of the collider phenomenology, such as event topology, lepton multiplicity, and missing energy, are highly dependent on the spectrum in the dark sector. Many of the suggested search strategies in [60], and past searches at the LEP and Tevatron will be relevant for the class of models considered here. In particular, if we restrict to a Higgs mass above the LEP bound, $m_h >$ 114 GeV [34], searches at the Tevatron for "dark photons" [61], next-to-minimal supersymmetric standard model pseudoscalars [62], trileptons [63], and same-sign dileptons [64] have the potential to constrain particular models and spectra.

To summarize, if moderate couplings to the portals on the order of the bottom quark Yukawa exist, there is an opportunity to probe the dark sector at high energy colliders via the decays of the SM Higgs boson. Looking forward, it seems useful to proceed in two directions due to the model dependence of the signatures. First, it would be worthwhile to study particular benchmark models containing some of the more complex signatures of multilepton, multijet final states with missing energy. Complementary to this, the parameter space of the more minimal models, particularly those with a single matter field, can be thoroughly explored.

V. DISCUSSION

Gauge symmetry provides a plausible origin for the discrete symmetries that can stabilize DM. In this paper we have surveyed a class of DM models based on an Abelian Z_N discrete gauge symmetry. We have investigated models with minimal field content, multifield Z_N models, and models with chiral matter. We have found that Z_N symmetries may lead to multiple stable DM candidates when N is not prime. The dark sector may couple to the SM through the kinetic mixing and Higgs portals, and we have given an outline of the basic phenomenology of this scenario, highlighting some novel direct detection and collider signatures. In particular, the SM Higgs boson may provide a portal through its decays to new states in the dark sector.

Our focus in this work has been on the gauge origin of the discrete symmetries and its potential impact on DM physics. There are a number of avenues for future work. Regarding indirect probes of these models, DM annihilation in the Galaxy and in the Sun can lead to novel astrophysical signatures. As is well known by now, a very light gauge boson provides a "dark force" [27] that can potentially explain the cosmic ray signatures of PAMELA [29] and the Fermi Gamma-ray Space Telescope [30]. In this context, there is a scope for the exploration of the low energy phenomenology of these models in fixed target [24,25] and e^+e^- experiments [22]. Also, DM annihilation in the Sun to light metastable mediators, such as dark Higgs or gauge bosons, can yield electromagnetic signatures in gamma-ray telescopes [65] such as the Fermi Gamma-ray Space Telescope [30]. In terms of cosmology, as recently pointed out in [66], GeV-scale states from a dark sector can help to resolve the lithium problem of big bang nucleosynthesis. It would be valuable to study these and other connections with cosmology and astrophysics for the DM models presented here.

In terms of model building, we have not attempted to address the electroweak hierarchy problem, or the additional naturalness issues that come with scalars in the dark sector. In this regard, it would be interesting to adapt these models to a supersymmetric framework along the lines of [59,67] or a warped-extra-dimensional framework [68]. This can dramatically alter the collider phenomenology, as now one may use the supersymmetric or Kaluza-Klein states as connectors to the dark sector.

We have focused on DM in a hidden sector, but one could also consider new U(1) gauge symmetries under which the SM is charged. There are a few anomaly-free global U(1) symmetries that one may gauge within the SM: $B - L$ (with right-handed neutrinos) or a difference in lepton numbers, e.g. $L_e - L_\mu$ [69]. However, more general U(1) symmetries can be gauged provided new fermion representations are added to cancel anomalies, and these could be used to obtain discrete gauge symmetries.

Stability (at least on cosmological time scales) is one of the few known properties of DM and hints at a new symmetry in the dark sector. Should experimental evidence for DM be found, new stabilization symmetries may be required to understand the data. In the meantime, exploring new possibilities for these symmetries may suggest new phenomena associated with DM and therefore new search strategies. Our investigation here illustrates this basic observation, as we have been led by symmetry considerations to predict new states, new interactions, and ultimately new signatures for direct detection and collider experiments.

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