Dark sector portal with vector-like leptons and flavor sequestering

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Abstract

We consider models with fermionic dark matter that transforms under a non-Abelian dark gauge group. Exotic, vector-like leptons that also transform under the dark gauge group can mix with standard model leptons after spontaneous symmetry breaking and serve as a portal between the dark and visible sectors. We show in an explicit, renormalizable model based on a dark SU(2) gauge group how this can lead to adequate dark matter annihilation to a standard model lepton flavor so that the correct relic density is obtained. We identify a discrete symmetry that allows mass mixing between the vector-like fermions and a single standard model lepton flavor, while preventing mixing between these fields and the remaining standard model leptons. This flavor sequestering avoids unwanted lepton-flavor-violating effects, substantially relaxing constraints on the mass scale of the vector-like states. We discuss aspects of the phenomenology of the model, including direct detection of the dark matter.

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I. INTRODUCTION

Although the literature on dark matter models is vast and diverse, the organizational structure of many of these models is similar. The visible sector includes all the fields normally associated with the minimal standard model; the dark sector consists of a collection of fields that communicate very weakly with the visible sector; the messenger or portal sector consists of those fields that allow for a weak coupling between the visible and dark sectors. In this paper, we are interested in a possible portal for dark matter models, specifically ones in which fermionic dark matter is charged under a dark gauge group. Our model will include vector-like fermions that are also in nontrivial representations of the dark gauge group but can mix with standard model fermions when the gauge symmetries of the theory are spontaneously broken. We identify a mechanism, based on symmetries, that we call "flavor sequestering" which allows this mixing to be non-negligible, while simultaneously suppressing unwanted flavor-changing processes. This mechanism is new to the literature; it can provide for vector-like fermion portal sectors that are lighter and more accessible experimentally than would otherwise be possible. For the purposes of illustration, we choose to study a theory with a non-Abelian dark gauge group, where an additional portal involving kinetic mixing of some dark gauge boson components with hypercharge is naturally suppressed. In models like the one we propose, where there are vector-like states charged both under the dark and hypercharge gauge groups, the kinetic mixing parameter in an Abelian theory would typically run below the Planck scale, leading to low-energy values that are not necessarily small; this makes a non-Abelian dark sector the natural setting for formulating our proposal. Scenarios in which multiple portals are relevant (for example, a vector-like fermion portal, a Higgs-portal, a higher-dimension-operator portal, *etc.*) are of course possible and more complicated; in the present work, however, we focus on the case where the vector-like fermion portal is dominant. Examples of non-Abelian dark matter models can be found in Refs. [1–7], [8–10] and [11]; we will not focus on models like those in Refs. [8–10] where a dark gauge boson is itself the dark matter. Our model is also very different from the models of Refs. [11] which involve unbroken non-Abelian dark gauge groups, either chosen to assure composite dark matter candidates in the cases where there is confinement, or dark radiation in the case where the dark gauge coupling is too small for confinement to be cosmologically relevant. In our proposal, mixing between the vector-like and standard model fermions will only be present when the non-Abelian dark gauge group is spontaneously broken.

Given these assumptions, we would like the vector-like fermion portal in our model to allow the dark gauge bosons to develop a small coupling to the visible sector, adequate enough to facilitate the annihilation of the dark matter for a successful thermal freeze-out, without running afoul of direct detection bounds. This can be arranged if the effective coupling between the dark and visible sectors does not appear at the same order in the dark matter annihilation and dark-matter-nucleon elastic scattering cross sections. To achieve this, we choose the quantum numbers of the vector-like states to allow mixing only with standard model leptons. The induced coupling of the dark gauge bosons would allow dark matter annihilation to leptons via tree-level diagrams, while diagrams involving quarks would be higher-order. One might wonder whether coupling the dark gauge bosons to standard model leptons directly might be a more economical alternative. However, proceeding in this way leads to significant model building complications. For example, if one tries to couple the dark gauge bosons to the standard model leptons directly, then the dark gauge bosons are potentially no longer "dark," unless their gauge coupling is taken to be very small. However, this choice suppresses the coupling of the dark gauge bosons to both the dark and visible sectors, making it ineffective as a channel for dark matter annihilation. Moreover, such direct couplings lead generically to chiral anomalies, which must be cancelled by additional states that are charged under both the dark and standard model gauge groups. There is no guarantee that the simplest Higgs field content of the dark and visible sectors will have the correct quantum numbers to provide Yukawa couplings for these additional states, so that additional Higgs representations may be required. Another potential problem is that charging standard model leptons under the new non-Abelian group may either restrict the form of the standard model lepton Yukawa matrices in unwanted ways, or forbid them entirely, unless a Higgs field charged under both the dark and standard model gauge groups is introduced. While the proliferation of fields implied by these considerations does not rise to the level of a no-go theorem, it does make the approach described a lot less appealing.

To avoid these complications, we assume that the non-Abelian dark gauge boson may couple to a vector-like state χ that can mix with standard model leptons after the gauge symmetries of the theory (both dark and visible) are spontaneously broken. We will refer to the χ states as heavy, vector-like leptons. If the dark gauge boson's coupling to dark matter is g_D, which may be substantial, then the induced coupling to the standard model lepton in the mass eigenstate basis will be proportional to θg_D where θ is a small mixing angle. Since the gauge boson couples directly to a vector-like state, anomalies are cancelled, and a mass term $-M_\chi \overline{\chi}\,\chi$ can be written down at tree-level. A case of particular phenomenological interest is where the vector-like sector is as light as possible. In this case, the mixing angle θ can be large enough so that the desired dark matter relic density is obtained entirely via dark matter annihilation to a standard model lepton-anti-lepton pair. This scenario would not be possible in a similar model without flavor sequestering, so we focus on this region of parameter space as the proof of principle that our flavor-sequestering idea can be incorporated in viable models. The range of M_χ is then determined by the requirement that the the mixing angle θ is large enough to produce the desired value of the dark matter relic density. In this paper, we will present an explicit and renormalizable model that illustrates this proposal. Our focus differs from that of Refs. [1–7], where the origin of higher-dimension operators connecting the dark and visible sectors was either unspecified, or assumed to arise from a sector whose flavor structure and phenomenology was not explicitly investigated.

The idea that a dark sector could communicate to the visible sector in any appreciable way through mixing between between vector-like leptons and their standard model counterparts would seem to conflict with the stringent lower bounds on the mass scale of heavy vector-like leptons that appear in the literature, which exceed 100 TeV [12]. Such stringent bounds, however, come from consideration of lepton-flavor-violating processes that emerge when the vector-like states mix with all three standard model lepton flavors. One expects such mixing to be present generically, and this would doom the approach that we have just outlined. In this paper, we show how a more favorable outcome can be achieved via discrete symmetries that allow us to suppress the unwanted mass mixing arbitrarily. In our model, vector-like leptons mix only with a single flavor of the standard model leptons, which in turn does not mix substantially with the remaining two flavors, thus avoiding problems with lepton flavor violation. We will show that the discrete symmetry used to achieve this flavor sequestering does not adversely affect the remaining flavor structure of the charged leptons or neutrino mass matrices. Phenomenological considerations place constraints on the mass spectrum of the flavor-sequestered vector-like lepton states that can be tested in direct collider searches.

Our paper is organized as follows. In the next section, we define the simplest model that illustrates a portal involving vector-like leptons and flavor sequestering. In Sec. 3, we show how the flavor structure of the theory can be achieved using a discrete symmetry, so that exclusive mixing with one standard model lepton generation is obtained and lepton-flavor- violating effects avoided. In Sec. 4 we demonstrate the viability of our example model by identifying the region of parameter space in which the correct dark matter relic density is obtained through annihilation to a standard model lepton-anti-lepton pair. We also consider the constraints from dark matter-nucleon elastic scattering, which follows from the suppressed kinetic mixing that is induced after the non-Abelian gauge group is spontaneously broken. In the final section, we summarize our conclusions.

II. THE MODEL

We consider the simplest non-Abelian dark gauge group, $SU(2)_D$. As stated earlier, we denote the heavy, vector-like leptons χ, and assume the quantum numbers

$$\chi_L \sim \chi_R \sim (\mathbf{2}, \mathbf{1}, \mathbf{1}, -1) \ , \tag{2.1}$$

where we indicate the representations of $SU(2)_D\times$ $SU(3)_C\times$ $SU(2)_W \times U(1)_Y$, in that order. In other words, these states are $SU(2)_D$ doublets, but have the same electroweak charges as right-handed leptons. We further assume the simplest assignment for the dark matter, *i.e.*, that it is a doublet under $SU(2)_D$. However, to avoid a Witten anomaly [13] there must be an even number of SU(2) fermion doublets, so we take

$$\psi_L \sim \psi_R \sim (\mathbf{2}, \mathbf{1}, \mathbf{1}, 0) \ . \tag{2.2}$$

Since the ψ fields are charged only under $SU(2)_D$, we can construct Dirac or Majorana mass terms, or both. We will assume Dirac mass terms, for simplicity, though it is easy to make this the only possibility by imposing additional discrete symmetries. For example, an unbroken Z_3 symmetry can forbid Majorana masses for ψ, and also serve as the symmetry which stabilizes the dark matter, which we identify henceforth as the lightest component of the ψ doublet.

We assume that the dark gauge symmetry is spontaneously broken by two $SU(2)_D$ Higgs field representations,

$$H_D \sim (\mathbf{2}, \mathbf{1}, \mathbf{1}, 0) \quad \text{and} \quad H_T \sim (\mathbf{3}, \mathbf{1}, \mathbf{1}, 0) \ . \tag{2.3}$$

We show at the end of this section that the Higgs potential has local minima consistent with the pattern of vacuum expectation values (vevs):

$$\langle H_D \rangle = \begin{pmatrix} v_{D1} \\ v_{D2} \end{pmatrix} \quad \text{and} \quad \langle H_T \rangle = \begin{pmatrix} v_T/2 & 0 \\ 0 & -v_T/2 \end{pmatrix} \quad . \tag{2.4}$$

If we decompose $H_T = H_T^a\,(\sigma^a/2)$, where the σ^a are Pauli matrices, then the H_T vev above corresponds to $\langle H_T^3 \rangle = v_T$ and $\langle H_T^a \rangle = 0$ for $a = 1, 2$. In fact, an arbitrary vev for H_T can always be rotated into the H_T^3 direction by an $\mathrm{SU}(2)_D$ transformation. With this choice, vevs in both components of H_D are expected, and one of those can be made real by a further $\mathrm{SU}(2)_D$ phase rotation. The fact that the remaining H_D vev in Eq. (2.4) is assumed real will be shown to be consistent with the minimization of a potential later.

We can now say something more concrete about the mass spectrum of the model. The relevant Lagrangian terms are $\mathcal{L} \supset \mathcal{L}_\psi + \mathcal{L}_{\chi e}$, where

$$\mathcal{L}_\psi = -M_\psi\,\overline{\psi}_L\,\psi_R + \lambda_s\,\overline{\psi}_L\,H_T\,\psi_R + \text{ h.c.} \quad , \tag{2.5}$$

and

$$\mathcal{L}_{\chi e} = -M_\chi\,\overline{\chi}_L\,\chi_R + \lambda_s'\,\overline{\chi}_L\,H_T\,\chi_R - y_1\,\overline{\chi}_L\,H_D\,e_R - y_2\,\overline{\chi}_L\,\widetilde{H}_D\,e_R - y_e\,\overline{L}_L\,H\,e_R + \text{ h.c.} \quad , \tag{2.6}$$

where $\widetilde{H}_D \equiv i\sigma^2 H_D^*$, and the final term is the usual standard model Yukawa coupling for a single lepton flavor. Eq. (2.6) assumes the existence of a symmetry that leads to exclusive mixing between any one standard model, right-handed charged lepton flavor (called e_R above) and the vector-like χ fields. We show how this flavor sequestering can be arranged by a discrete symmetry in Sec. III. The first terms in Eqs. (2.5) and (2.6) provide a common mass for each component of the given doublet, while the second terms lead to mass splittings proportional to the vev v_T. The third and fourth terms in Eq. (2.6) allow mixing between the standard model lepton e_R and the χ fields, since the coupling to the dark doublet Higgs field H_D allows for the formation of an $\mathrm{SU}(2)_D$ singlet. The final term leads to an e mass when the standard model Higgs field develops a vacuum expectation value $\langle H \rangle = (0, v/\sqrt{2})$, with $v = 246$ GeV. Defining the column vector $\Upsilon \equiv (e,\,\chi^{(1)},\,\chi^{(2)})^T$, which displays the two components of the χ doublet, we may write the mass matrix that is produced after spontaneous symmetry breaking by

$$\mathcal{L}_{mass}^{\chi e} = -\overline{\Upsilon}_L M\,\Upsilon_R + \text{ h.c.} \quad , \tag{2.7}$$

where

$$M = \begin{pmatrix} \frac{h_e v}{\sqrt{2}} & 0 & 0 \\ \frac{(y_1 v_{1D} + y_2 v_{2D})}{\sqrt{2}} & M_\chi - \frac{\lambda'_s v_T}{2} & 0 \\ \frac{(y_1 v_{2D} - y_2 v_{1D})}{\sqrt{2}} & 0 & M_\chi + \frac{\lambda'_s v_T}{2} \end{pmatrix} \equiv \begin{pmatrix} m_0 & 0 & 0 \\ m_1 & M_1 & 0 \\ m_2 & 0 & M_2 \end{pmatrix} , \tag{2.8}$$

where the second form is a convenient parametrization. This matrix can be diagonalized by a bi-unitary transformation, $M = U_L \, M^{diag} \, U_R^\dagger$. While this diagonalization can be done numerically, there are certain limits that are relevant to us in which simple results can be obtained. In particular, when M_1, $M_2 >> m_1$, $m_2 >> m_0$, we find that the largest mixing angles, which occur in U_R, are given by

$$U_R = \begin{pmatrix} 1 - \frac{1}{2}\left(\frac{m_1^2}{M_1^2} + \frac{m_2^2}{M_2^2}\right) & m_1/M_1 & m_2/M_2 \\ -m_1/M_1 & 1 - \frac{1}{2}\frac{m_1^2}{M_1^2} & -\frac{M_1}{M_2}\frac{m_1 m_2}{M_1^2 - M_2^2} \\ -m_2/M_2 & \frac{M_2}{M_1}\frac{m_1 m_2}{M_1^2 - M_2^2} & 1 - \frac{1}{2}\frac{m_2^2}{M_2^2} \end{pmatrix} + \cdots , \tag{2.9}$$

where the $\cdots$ represent terms that are cubic order or higher in m_i/M_j. For this case, we can now find the leading coupling of the dark gauge fields $A^a_{D\mu}$ to the mass eigenstate fields. In the gauge basis, the coupling to Υ_R can be written

$$\mathcal{L} = i\overline{\Upsilon}_R \gamma^\mu (\partial_\mu - i g_D A^a_{D\mu} \mathcal{T}^a) \Upsilon_R + \cdots , \tag{2.10}$$

where

$$\mathcal{T}^a = \left(\begin{array}{c|c} 0 & 0 \\ \hline 0 & T^a \end{array}\right) , \tag{2.11}$$

and $T^a = \sigma^a/2$, $a = 1, \ldots, 3$, are the generators of SU(2). The zero in the 1-1 element reflects the fact that the standard model lepton is not charged under the dark gauge group. In the mass eigenstate basis, the couplings of the a^{th} dark gauge boson are therefore proportional to $U_R^\dagger \mathcal{T}^a U_R$. In the same approximation as Eq. (2.9), these matrices are given by

$$U_R^\dagger \mathcal{T}^a U_R = \left[\begin{pmatrix} \frac{m_1 m_2}{M_1 M_2} & -\frac{m_2}{2M_2} & -\frac{m_1}{2M_1} \\ -\frac{m_2}{2M_2} & 0 & \frac{1}{2} \\ -\frac{m_1}{2M_1} & \frac{1}{2} & 0 \end{pmatrix} , \begin{pmatrix} 0 & -\frac{i\,m_2}{2M_2} & \frac{i\,m_1}{2M_1} \\ \frac{i\,m_2}{2M_2} & 0 & -\frac{i}{2} \\ -\frac{i\,m_1}{2M_1} & \frac{i}{2} & 0 \end{pmatrix} , \begin{pmatrix} \frac{m_1^2}{2M_1^2} - \frac{m_2^2}{2M_2^2} & -\frac{m_1}{2M_1} & \frac{m_2}{2M_2} \\ -\frac{m_1}{2M_1} & \frac{1}{2} & 0 \\ \frac{m_2}{2M_2} & 0 & -\frac{1}{2} \end{pmatrix} \right] , \tag{2.12}$$

where we only show results to linear order in m_i/M_j, with the exception of the 1-1 entries, because of their relevance to our subsequent discussion. For example, for the lightest dark

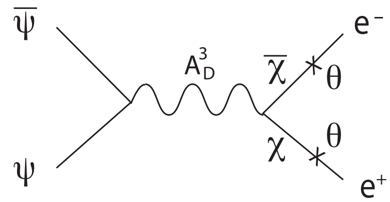

FIG. 1: Qualitative picture of dark matter annihilation to a charged lepton-anti-lepton pair, as discussed in the text. The insertions labelled by θ represent mass mixing.

gauge boson, A_D^3, the coupling to $e^+\, e^-$ is given by

$$g_D \overline{\Upsilon}_R \gamma^\mu A^3_{D\mu} (U_R^\dagger \mathcal{T}^3 U_R) \Upsilon_R = \frac{g_D}{2} \left(\frac{m_1^2}{M_1^2} - \frac{m_2^2}{M_2^2} \right) \overline{e}_R \gamma^\mu A^3_{D\mu} e_R + \cdots \tag{2.13}$$

which provides the A_D^3 gauge boson with a decay channel (since we assume its mass is greater that $2\, m_e$) and allows for dark matter annihilation to a charged standard model lepton-anti-lepton pair. For later convenience, we define

$$\theta^2 \equiv g_D \left(\frac{m_1^2}{M_1^2} - \frac{m_2^2}{M_2^2} \right) \quad . \tag{2.14}$$

We illustrate the qualitative idea in Fig. 1 that the dark matter annihilation to a charged standard model lepton-anti-lepton pair emerges from mixing that affects two of the external legs.

We note that in the case where m_0 is comparable to m_1 and m_2 we find via numerical diagonalization that our expression U_R in Eq. (2.9) still provides an accurate approximation. Moreover, we can prove that m_0 appears only as a higher-order correction to θ, as defined in Eq. (2.14), the quantity that is most relevant to our phenomenological discussion later. The argument is as follows: if m_1 or m_2 where to vanish, then U_R must become the identity. This implies that any corrections to the 1-2, 1-3, 2-1 and 3-1 entries of U_R that are proportional to m_0 must come at no lower order than $m_0 m_{1,2}/M_{1,2}^2$. This potential contribution is nonetheless higher-order than the values shown for these entries in Eq. (2.9). It is also the case that the 1-1 entry of $U_R^\dagger \mathcal{T}^3 U_R$, from which θ is extracted, depends only on these four entries. Hence, the value of θ, which controls the induced coupling of A_D^3 to the chosen standard model lepton flavor, remains unaffected at leading order.

Eq. (2.12) indicates that all states other than the lightest ψ mass eigenstate have available decay channels that ultimately lead to standard model particles. Since the free parameter space of our model is substantial, for definiteness we assume henceforth the following about the spectrum:

- Due to the triplet vev, $\psi^{(1)}$ and $A_D^{(3)}$ are the lightest states of the dark sector, while $\psi^{(2)}$, $A_D^{(1)}$ and $A_D^{(2)}$ are substantially heavier. We will consider the case where the lighter dark sector states are in the $\mathcal{O}(1) - \mathcal{O}(100)$ GeV range, with the restriction that $m_{\psi^{(1)}} \leq m_{A_D^{(3)}}$, so that the dominant dark matter annihilation channel proceeds through the vector-like lepton portal (see Sec. IV).

- We assume that the vector-like leptons χ have masses above $M_Z/2$, so that the Z width is unaffected. Note that more substantial collider bounds apply when vector-like leptons are either in weak doublets, or are long-lived [16], neither of which applies in the present case.

With these assumptions, let us first sketch out the decay modes when the standard model lepton flavor involved is the electron: the coupling matrices $U_R^\dagger \mathcal{T}^a U_R$, for $a = 1$ and $a = 3$ allow decays of A_D^1 and A_D^3 directly to $e^+ e^-$; the same is not true for $a = 2$, but the A_D^2 boson does couple to the two different ψ mass eigenstates, which we will call $\psi^{(1)}$ (the lighter, dark matter component) and $\psi^{(2)}$ (the heavier). The eigenstate $\psi^{(2)}$ can decay to dark matter $\psi^{(1)}$ plus $e^+ e^-$ via A_D^1 exchange. Hence A_D^2 can decay to two dark matter particles and an $e^+ e^-$ pair, whether or not $\psi^{(1)}$ is on shell. Due to the $\chi A_D^3 e$ couplings in $U_R^\dagger \mathcal{T}^3 U_R$, both χ mass eigenstates can decay to a same-sign e plus an $e^+ e^-$ pair via A_D^3 exchange. Finally, the exotic Higgs fields H_D and H_T couple to fermion pairs via their Yukawa couplings. Since we have already established that those fermions couple ultimately to either e's or $\psi^{(1)}$'s, our claim is established. Note that if the standard model fermion is either μ or τ, nothing above is changed if $M_{A_D^3} > 2m_\mu$ or $2m_\tau$; otherwise, decays to lighter charged leptons plus neutrinos can still occur with the μ's or τ's off shell.

Since the χ and e_R have identical electroweak quantum numbers, there is no effect on the coupling of the Z boson to e_R in the mass eigenstate basis. However, χ and e_L couple differently to the electroweak gauge bosons, and diagonalization of Eq. (2.8) also involves a left-handed rotation matrix U_L which differs from the identity. Fortunately, the left-handed mixing angles are much smaller than those in Eq. (2.9) so that this does not present any phenomenological difficulties. For example, the fractional shift in the standard model Ze_Le_L vertex is of $\mathcal{O}(\frac{m_0 m_1}{M_1^2} \frac{m_0 m_2}{M_2^2})$, which is negligible given the spectrum we assume in Sec. IV. We also may take the mostly χ mass eigenstates to be heavy enough so that rare Z decays to $\chi\, e$ are kinematically forbidden.

Finally, let us return to the issue of the spontaneous breaking of the dark gauge symmetry. In the effective theory well below the electroweak scale, the most general renormalizable potential involving the dark Higgs fields is given by

$$\begin{aligned} V(H_D, H_T) = & -m_D^2 H_D^\dagger H_D - m_T^2\, \mathrm{tr}(H_T H_T) + \lambda_1 (H_D^\dagger H_D)^2 + \lambda_2 \big[\, \mathrm{tr}(H_T H_T)\big]^2 \\ & + \lambda_3 H_D^\dagger H_T H_T H_D + \mu_1 H_D^\dagger H_T H_D + \Big(\mu_2 H_D^\dagger H_T \widetilde{H}_D + \mathrm{h.c.}\Big)\,, \end{aligned} \tag{2.15}$$

where we have used the fact that $H_T^\dagger = H_T$. We assume the potential does not violate CP, so that all the couplings are real. Further, we require at least one of $(-m_D^2,\ -m_T^2)$ to be negative so that the H_D and H_T fields may develop non-zero vevs. It should be noted that there are other terms involving the Higgs fields that could be added to the potential, such as $\widetilde{H}_D^\dagger \widetilde{H}_D$, $\mathrm{tr}(H_T^4)$, $\widetilde{H}_D^\dagger H_T \widetilde{H}_D$, $H_D^\dagger H_D\, \mathrm{tr}(H_T H_T)$, but these are not linearly independent of the terms included in Eq. (2.15) and so have been omitted.

The Higgs doublet assumes the standard real-field parametrization,

$$H_D = \frac{1}{\sqrt{2}} \begin{pmatrix} \phi_1 + i\phi_2 \\ \phi_3 + i\phi_4 \end{pmatrix}, \tag{2.16}$$

while the Higgs triplet can be represented by a 2×2 matrix of real fields H_1, H_2 and H_3,

$$H_T = H^a \frac{\sigma^a}{2} = \frac{1}{2} \begin{pmatrix} H_3 & H_1 - iH_2 \\ H_1 + iH_2 & -H_3 \end{pmatrix}. \tag{2.17}$$

The normalization assures canonical kinetic terms. We proceed to show that there exists a stable, local minimum of the potential for the pattern of vacuum expectation values described in Eq. (2.4). One approach to studying the potential is to fix all the parameters and search for minima, using standard steepest descent algorithms. However the downside to this approach is that one may then have to repeatedly discard local minima that do not provide the pattern of vevs desired for the model. So instead, we will fix the vevs and work backwards, showing that an extremum exists that is also a local minimum for a fixed set of parameters.

The extremization of Eq. (2.15) with the fields set to the vevs shown in Eq. (2.4) provides the following nontrivial, linearly independent constraints:

$$
\begin{aligned}
-m_D^2 v_{D1} + \lambda_1 v_{D1}^3 + \lambda_1 v_{D1} v_{D2}^2 + \mu_2 v_{D2} v_T + \frac{1}{4} v_{D1} v_T (\lambda_3 v_T + 2\mu_1) &= 0 \\
\frac{1}{2}\left(-\mu_2 v_{D1}^2 + \mu_1 v_{D1} v_{D2} + \mu_2 v_{D2}^2\right) &= 0 \\
-m_T^2 v_T + \frac{1}{4}\lambda_3 v_T \left(v_{D1}^2 + v_{D2}^2\right) + \frac{1}{4}\mu_1 (v_{D1}^2 - v_{D2}^2) + \mu_2 v_{D1} v_{D2} + \lambda_2 v_T^3 &= 0.
\end{aligned}
\tag{2.18}
$$

For the purpose of numerical evaluation we work here in units where $\mu_1 = 1$. For fixed choices of the vevs and the couplings $\lambda_{1,2,3}$, we may then determine m_D, m_T and μ_2. To determine whether the extremum is a minimum, maximum or saddle point, we need to examine the eigenvalues of the mass squared matrix (the second derivative matrix with all fields set to their vevs and with the solutions for m_D, m_T and μ_2 corresponding to the extremum). Since $SU(2)_D$ is spontaneously broken to nothing, we expect three Goldstone bosons, one for each broken SU(2) generator, according to Goldstone's theorem. Thus we would expect three of the eigenvalues to be zero, corresponding to the massless degrees of freedom that are "eaten" by the dark gauge bosons. The remaining eigenvalues must be positive for the extremum to be a local minimum. For example, let us set $v_T = v_{D1} = v_{D2}/2 = \lambda_{1,2,3} = \mu_1$ (here we require $v_{D1} \neq v_{D2}$ for a solution to exist). Then we find $m_D^2 = 53/12$, $m_T^2 = 1/6$ and $\mu_2 = -2/3$. The corresponding mass squared eigenvalues are $\{0, 0, 0, 3.75, 3.75, 4, 10\}$, in units of μ_1^2, thus confirming that we are at a local minimum of the potential. This provides an existence proof that local minima exist in which the pattern of vevs shown in Eq. (2.4) is obtained. It is not difficult to find similar solutions for other choices of v_{D1}, v_{D2} and v_T.

The SU$(2)_D$ breaking vevs affect the χ-e mass spectrum via Eq. (2.8); the triplet vev also splits the ψ mass eigenstates

$$
m_{\psi^{(1)}} = M_\psi - \frac{1}{2}\lambda_s v_T, \quad m_{\psi^{(2)}} = M_\psi + \frac{1}{2}\lambda_s v_T \tag{2.19}
$$

for $\psi_{L,R} = \left(\psi^{(1)},\ \psi^{(2)}\right)^T_{L,R}$. The gauge field spectrum is obtained from the kinetic terms for H_D and H_T,

$$
\mathcal{L}_{kin}(H_D, H_T) = (D_\mu H_D)^\dagger (D^\mu H_D) + \mathrm{tr}\left[(D_\mu H_T)^\dagger (D^\mu H_T)\right], \tag{2.20}
$$

where $D_\mu H_D = \partial_\mu H_D - i g_D A^a_{D\mu} \frac{\sigma^a}{2} H_D$ and $D_\mu H_T = \partial_\mu H_T - i g_D \frac{\sigma^a}{2} A^a_{D\mu} H_T + i g_D A^a_{D\mu} H_T \frac{\sigma^a}{2}$. Following symmetry breaking the gauge bosons develop masses

$$
m^2_{A^1_D} = m^2_{A^2_D} = \frac{g_D^2}{4}(v_{D1}^2 + v_{D_2}^2 + 4\, v_T^2), \quad m^2_{A^3_D} = \frac{g_D^2}{4}(v_{D1}^2 + v_{D_2}^2). \tag{2.21}
$$

In splitting the ψ and A_D multiplet masses, the triplet vev leads to a simple low-energy effective theory consisting of the dark matter $\psi^{(1)}$ (we assume $\lambda_s > 0$) and the mediator A_D^3, which has small induced couplings to a right-handed standard model lepton flavor. This effective theory is relevant below the masses of the heavy vector-like leptons, $\psi^{(2)}$ and the $A_D^{1,2}$ bosons, which we will associate with a common scale, for simplicity. In addition, we will see that the triplet vev leads to induced couplings of the dark matter to quarks via kinetic mixing, which will lead to avenues for direct detection. We discuss the phenomenology of this scenario in Sec. IV.

III. FLAVOR SEQUESTERING

In this section, we show that it is possible to allow for non-negligible mixing between one flavor of the standard model leptons and the heavier, vector-like leptons, while suppressing the mixing with the other standard-model flavors, so that bounds on lepton-flavor-violating processes become irrelevant. In the discussion below, we refer to that one flavor as the electron e, though the approach described applies equally well if the chosen flavor were μ or τ. Let us consider the structure of the standard model Yukawa matrices first, and then introduce couplings to the vector-like states.

We represent the three generation of standard model lepton doublets by L_{iL} and the right-handed charged leptons by E_{iR}, for $i = 1, \ldots, 3$, We imagine that the Yukawa couplings are determined by a flavor symmetry of the form $Z_N \times G_F$. Our interest is in the effect of the Z_N factor, while we do not commit to any specific G_F. We aim to show that the restrictions that follow from the Z_N symmetry are sufficient to suppress the flavor mixing effects that we would like to avoid, while remaining compatible with a variety of possible flavor models that may determine the remaining, detailed structure of the Yukawa matrices.

We represent an element of Z_N by ω^j, for $j = 1, \ldots N$, where $\omega^N \equiv 1$. We assign the following transformation properties to the L and E fields, representing them here as column vectors:

$$L_L \to \Omega \, L_L \quad \text{and} \quad E_R \to \Omega \, E_R \ , \tag{3.1}$$

where

$$\Omega = \begin{pmatrix} 1 & 0 & 0 \\ 0 & \omega^{-n} & 0 \\ 0 & 0 & \omega^{n} \end{pmatrix} \quad . \tag{3.2}$$

Note that $\omega^{-n} \equiv \omega^{N-n}$. Assuming that the standard model Higgs doublet is unaffected by the Z_N symmetry, the transformation properties of the charged-lepton Yukawa matrix entries that lead to invariant couplings are summarized by

$$Y_E \sim \left(\begin{array}{c|cc} 1 & \omega^{n} & \omega^{-n} \\ \hline \omega^{-n} & 1 & \omega^{-2n} \\ \omega^{n} & \omega^{2n} & 1 \end{array}\right) \quad , \tag{3.3}$$

where the transformation property of, for example, the 1-2 entry is understood to be $Y_E^{12} \to \omega^n \, Y_E^{12}$, and so on. We will choose $N = 2\,n$ so that the entire two-by-two block on the lower right is unconstrained by the Z_N symmetry, the least restrictive possibility that meets our needs[1]. The amount by which the electron mass eigenstate is affected by the second and third generation fields, however, is entirely controlled by the size of n, once Z_N breaking fields are introduced, as we discuss later.

A symmetry affecting the left-handed charged leptons also affects the left-handed neutrinos, so we must verify that neutrino phenomenology is not adversely affected. For example, if we had imposed a Z_2 symmetry, with $n = 1$, and required it to remain exactly unbroken, we can also completely eliminate mixing between the first generation charged leptons and those of the second and third generations. However, if we then introduce three generations of right-handed neutrinos N_i, for $i = 1, \ldots, 3$, one can show that there are no Z_2 charge assignments for the N fields that leads to the correct neutrino mass squared differences and mixing angles, assuming the light mass eigenstates follow from the see-saw mechanism. However, more favorable results may be obtained when the Z_N symmetry is spontaneously broken. Here, we assume the same transformation for all three right-handed neutrino fields:

$$N_R \to \omega^p \, N_R \quad , \tag{3.4}$$

[1] This choice is also compatible with G_F having a non-Abelian component in which two flavors of standard model leptons transform as a doublet. However, it is sufficient (and simplest) for present purposes to imagine that G_F has only Abelian factors.

where p is an integer. Defining the Dirac neutrino mass via $\mathcal{L} \supset \overline{L}_L \widetilde{H} Y_{LR} N_R + \text{h.c.}$, the transformation properties of the Yukawa coupling is given by

$$Y_{LR} \sim \left(\begin{array}{c|cc} \omega^{-p} & \omega^{-p} & \omega^{-p} \\ \hline \omega^{-n-p} & \omega^{-n-p} & \omega^{-n-p} \\ \omega^{n-p} & \omega^{n-p} & \omega^{n-p} \end{array} \right) \quad . \tag{3.5}$$

For the choice $n = 2\,p$, or equivalently $N = 2\,n = 4\,p$, we may use the fact that $\omega^{-n-p} \equiv \omega^p$ and $\omega^{n-p} \equiv \omega^p$ to write

$$Y_{LR} \sim \left(\begin{array}{c|cc} \omega^{-p} & \omega^{-p} & \omega^{-p} \\ \hline \omega^{p} & \omega^{p} & \omega^{p} \\ \omega^{p} & \omega^{p} & \omega^{p} \end{array} \right) \quad . \tag{3.6}$$

The significance of this form is clear if we assume that there is a flavon field ρ with the Z_N transformation property

$$\rho \to \omega\, \rho \ \ , \tag{3.7}$$

and a vacuum expectation value such that $\langle \rho \rangle / M \equiv \epsilon$ is a small parameter. Here M is the flavor scale, which is the ultraviolet cut off of the effective theory. Then all the entries of Y_{LR} are non-vanishing, and proportional to either $(\rho/M)^p$ or to $(\rho^*/M)^p$. Hence, we may write

$$Y_{LR} = \epsilon^p\, \widetilde{Y}_{LR} \ \ , \tag{3.8}$$

where $\widetilde{Y}_{LR}$ is a three-by-three matrix that is thus far arbitrary. Following a similar argument, we define the right-handed neutrino Majorana mass matrix by the Lagrangian term $\overline{N^c}_R M_{RR} N_R$, and see immediately that

$$M_{RR} \to \omega^{-2p} M_{RR} \ \ . \tag{3.9}$$

Again, this is consistent with the transformation property of $(\rho^*/M)^{2p}$, so we may write

$$M_{RR} = \epsilon^{2p} \widetilde{M}_{RR} \ \ , \tag{3.10}$$

where $\widetilde{M}_{RR}$ is a three-by-three Majorana mass matrix that is also arbitrary thus far. With $\widetilde{Y}_{LR}$ and $\widetilde{M}_{RR}$ arbitrary, it is possible to obtain any desired neutrino phenomenology, which demonstrates that the Z_N symmetry does not lead to unwanted phenomenological restrictions. Theories that predict the detailed structure of $\widetilde{Y}_{LR}$ and $\widetilde{M}_{RR}$ by the breaking of an additional symmetry G_F are compatible with this framework. Note that the overall powers of ϵ in Eqs. (3.8) and (3.10) scale out of the see-saw formula which determines the Majorana mass matrix for the three light neutrino mass eigenstates

$$M_{LL} = M_{LR}\, M_{RR}^{-1}\, M_{LR}^T, \tag{3.11}$$

where $M_{LR} = (v/\sqrt{2})\, Y_{LR}$. The effect of the Z_N symmetry on the form of the charged lepton Yukawa matrix is to impose the form

$$Y_E \sim \left(\begin{array}{c|cc} y_{11} & \epsilon^n \tilde{y}_{12} & \epsilon^n \tilde{y}_{13} \\ \hline \epsilon^n \tilde{y}_{21} & y_{22} & y_{23} \\ \epsilon^n \tilde{y}_{31} & y_{32} & y_{33} \end{array}\right) \;. \tag{3.12}$$

For ϵ sufficiently small, or n sufficiently large, or both, we can make Y_E as close to block diagonal as we like.

Now we include the vector-like state χ with the same electroweak quantum numbers as a right-handed electron, but charged also under a dark gauge group. Yukawa couplings involving $\overline{\chi}_L e_R$ and a dark Higgs field are unaffected by the Z_N symmetry, while those involving $\overline{\chi}_L \mu_R$ or $\overline{\chi}_L \tau_R$ transform by $\omega^{\pm n}$. These potential sources of unwanted mixing that may emerge after the dark gauge symmetry is spontaneously broken are therefore highly suppressed by the same factors of ϵ^n that appear in the unwanted entries in Y_E. We conclude that it is possible to make the χ, e, μ, τ mass matrix as block diagonal as desired, by suitable choice of ϵ^n, such that χ mixes substantially only with e, or any one desired lepton flavor, by a similar construction[2].

The question of which lepton flavor is selected to mix with the heavier, vector-like states impacts the phenomenology of the dark gauge bosons. For example, if the mixing only involves the τ lepton, then bounds on the A_D^a from searches for s-channel resonances in low-energy e^+e^- collisions, or from indirect processes like the electron or muon $g-2$ would be irrelevant. The phenomenology in the case where the mixing involves either a first or

[2] We also note that this result is not linked in any fundamental way to our initial choice to study a non-Abelian dark gauge group. The present approach would be equally effective if the χ mass mixing terms were generated after the spontaneous breaking of a dark Abelian gauge symmetry. However, as noted earlier, Abelian theories would generically have kinetic mixing with hypercharge at tree-level and one-loop running of the mixing parameter below the Planck scale induced by the presence of the vector-like lepton states. The flavor sequestering mechanism could be applied in Abelian dark sector models provided that an additional mechanism is specified that adequately suppresses these kinetic mixing effects.

second generation lepton would lead to more meaningful constraints, but one that would depend on other assumptions about the spectrum, for example if A_D^3 decays visibly or invisibly, which depends on the dark matter mass. In the following section, we will assume the least constrained possibility, that the χ's mix with the τ, and consider the wide range of phenomenological issues associated with the other two possibilities in separate work [14]. This has the appealing aesthetic feature that the flavor symmetry distinguishes the third generation from the other two, an idea that has appeared in many other contexts in the literature on the flavor structure of the standard model [15].

IV. PHENOMENOLOGY

To confirm the viability of our flavor-sequestered model, we wish to show that it can achieve the correct dark matter relic density. We will not do a complete study of the model's parameter space, but focus on a region that is unique to the flavor-sequestered scenario, namely where the vector-like states are light enough so that sufficient dark matter annihilation is achieved to standard model lepton-anti-lepton pairs, even in the absence of other annihilation channels. We then comment on direct detection via the suppressed kinetic mixing effects that emerge when the gauge symmetries are broken.

A. Relic Density

The scattering amplitude for s-channel dark matter annihilation into standard model particles depicted in Fig. 1, with e replaced by τ, is given by

$$\mathcal{M}(\psi^{(1)}\overline{\psi^{(1)}} \to \tau^+\tau^-) = \frac{i g_D^2 \theta^2}{4\big(q^2 - m_{A_D^3}^2 + i m_{A_D^3}\Gamma^D\big)}\ \overline{v}(p')\gamma^\mu u(p)\ \overline{u}(k)\gamma_\mu v(k') \tag{4.1}$$

where p (p') is the momentum of the incoming dark matter fermion (anti-fermion), k (k') is the the momentum of the outgoing τ^- (τ^+) and $q = p + p'$ is the momentum flowing through the A_D^3 propagator. As discussed in Sec. II, the lightest gauge boson A_D^3 couples to the vector-like states $\chi^{(1)}$ and $\chi^{(2)}$, which then mix with a standard model lepton flavor (chosen here as τ) after spontaneous symmetry breaking. This results in a factor of θ^2, defined in Eq. (2.14), in the scattering amplitude.

Our numerical results for dark matter annihilation depend on assumptions about the dark particle mass spectrum and couplings. We assume the picture described earlier, where the lightest states consist of $\psi^{(1)}$ and A_D^3, and decays of A_D^3 to any of the heavier exotic states are not kinematically allowed. For the mass range studied in this section, A_D^3 can decay to $\tau^+\tau^-$, and possibly also $\psi^{(1)}\overline{\psi^{(1)}}$, depending on the dark matter mass. Consequently, the total decay width of the dark gauge boson appearing in the propagator is given by

$$\Gamma^D = \Gamma\left(A_D^3 \to \tau^+\tau^-\right) + \Theta\left(m_{A_D^3} - 2\, m_{\psi^{(1)}}\right)\, \Gamma\left(A_D^3 \to \psi^{(1)}\overline{\psi^{(1)}}\right) \tag{4.2}$$

where Θ is a step function, *i.e.*, $\Theta(x) = 1$ if $x \geq 0$ and $\Theta(x) = 0$ if $x < 0$, and

$$\Gamma\left(A_D^3 \to \tau^+\tau^-\right) = \frac{1}{48\pi} g_D^2 m_{A_D^3} \theta^4 \left(1 + \frac{2\, m_\tau^2}{m_{A_D^3}^2}\right)\left(1 - \frac{4\, m_\tau^2}{m_{A_D^3}^2}\right)^{1/2}, \tag{4.3}$$

$$\Gamma\left(A_D^3 \to \psi^{(1)}\overline{\psi^{(1)}}\right) = \frac{1}{48\pi} g_D^2 m_{A_D^3} \left(1 + \frac{2\, m_{\psi^{(1)}}^2}{m_{A_D^3}^2}\right)\left(1 - \frac{4\, m_{\psi^{(1)}}^2}{m_{A_D^3}^2}\right)^{1/2}. \tag{4.4}$$

Since the mean dark matter velocity is typically around 220 km/s [16], we work in the non-relativistic limit where $E_{\psi^{(1)}} \approx m_{\psi^{(1)}}$. We then find the thermally averaged annihilation cross section times velocity

$$\langle \sigma_A v \rangle = \frac{g_D^4 \theta^4}{32\pi} \frac{2m_{\psi^{(1)}}^2 + m_\tau^2}{(4m_{\psi^{(1)}}^2 - m_{A_D^3}^2)^2 + m_{A_D^3}^2 \Gamma_D^2} \left(1 - \frac{m_\tau^2}{m_{\psi^{(1)}}^2}\right)^{1/2}. \tag{4.5}$$

Using this we calculate the freeze-out temperature T_F and the dark matter relic density by standard methods [17]. Dark matter freeze out occurs when the interaction probability per unit time $\Gamma_{\psi^{(1)}}$, equals the expansion rate of the universe, H, *i.e.*,

$$\left.\frac{\Gamma_{\psi^{(1)}}}{H}\right|_{T=T_F} = \left.\frac{n_{EQ}^{\psi^{(1)}} \langle \sigma_A v \rangle}{H}\right|_{T=T_F} \simeq 1. \tag{4.6}$$

Here $n_{EQ}^{\psi^{(1)}}$ is the equilibrium number density of the dark matter particle, given by

$$n_{EQ}^{\psi^{(1)}} = 2\left(\frac{m_{\psi^{(1)}} T}{2\pi}\right)^{3/2} e^{-m_{\psi^{(1)}}/T}. \tag{4.7}$$

Freeze-out occurs during the radiation-dominated epoch in which case

$$H = 1.66\, g_*^{1/2} T^2 / M_{pl}, \tag{4.8}$$

where $M_{pl} = 1.22 \times 10^{19}$ GeV is the Planck mass and $g_*(T)$ the number of relativistic degrees of freedom at temperature T,

$$g_*(T) = \sum_{i=bosons} g_i \left(\frac{T_i}{T}\right)^4 + \frac{7}{8} \sum_{i=fermions} g_i \left(\frac{T_i}{T}\right)^4. \tag{4.9}$$

Finally the dark matter relic density is given by

$$\Omega_D h^2 = \frac{2 \cdot (1.07 \times 10^9 \text{ GeV}^{-1})\, x_F}{\sqrt{g_*(T_F)} M_{Pl} \langle \sigma_A v \rangle}. \tag{4.10}$$

We define $x_F \equiv m_{\psi^{(1)}}/T_F$ where T_F is obtained by solving Eq. (4.6). The factor of 2 is included because we are accounting for the density of dark matter particles and antiparticles. We require Eq. (4.10) to reproduce the WMAP result 0.1186 ± 0.0020 [16] within two standard deviations.

To display our results, we fix $m_{A_D^3}$ and θ and find the regions of the g_D-$m_{\psi^{(1)}}$ plane in which the desired dark matter relic density is obtained. We assume that the mixing angle remains small ($\theta < 1$) but not so small that a satisfactory dark matter annihilation cross section cannot be obtained. So that the dark gauge coupling remains perturbative, we assume $\alpha_D/(4\pi) < 1/3$ or equivalently $g_D < 4\pi/\sqrt{3} \sim 7.25$; one-loop corrections become comparable to tree-level amplitudes when $\alpha/(4\pi) \approx 1$, so one-third of this value is a reasonable upper limit on the dark coupling constant. For the purposes of determining g_*, we assume all exotic mass eigenstates other than $\psi^{(1)}$ and A_D^3, are at $m_Z = 91.1876$ GeV. With this choice, the Z boson cannot decay into $\chi\overline{\chi}$ or $\chi\tau$, which could lead to an unacceptable broadening of the precisely measured Z boson width [16].

Fig. 2 shows the regions of the g_D-$m_{\psi^{(1)}}$ plane in which the dark matter relic density is within two standard deviations of the WMAP result 0.1186 ± 0.0020 [16], for fixed choices of $m_{A_D^3}$ and θ. We have intentionally centered the plots around the point of resonance annihilation $m_{\psi^{(1)}} = m_{A_D^3}/2$ where the cross section is largest. For small values of g_D at fixed θ, some tuning is required to achieve a large enough annihilation cross section. However, Fig. 2 indicates that we can have larger, perturbative values of g_D without requiring that we sit unnaturally close to the resonance. As θ is made progressively smaller, however, more tuning is required. This is indicated by the narrowing range in $m_{\psi^{(1)}}$ for each solution in which g_D is also perturbative.

Of course, the values of θ that are indicated in Fig 2 are related to choices for the masses and coupling in the model, such that $\theta^2 = g_D \left(\frac{m_1^2}{M_1^2} - \frac{m_2^2}{M_2^2} \right)$, where the m_i and M_i were defined in Eq. (2.8). It is not hard to verify that the values of θ shown in Fig. 2 can be achieved given the assumptions that went into the making of the plots. For example, in the $m_{A_D^3} = 10$ GeV plot, consider the point where $g_D \approx 1$ and $m_{\psi^{(1)}} \approx 8.5$ GeV, on the $\theta = 0.1$ band. Given our earlier assumption in computing g_* that the heavier exotic states are at m_Z, one can check

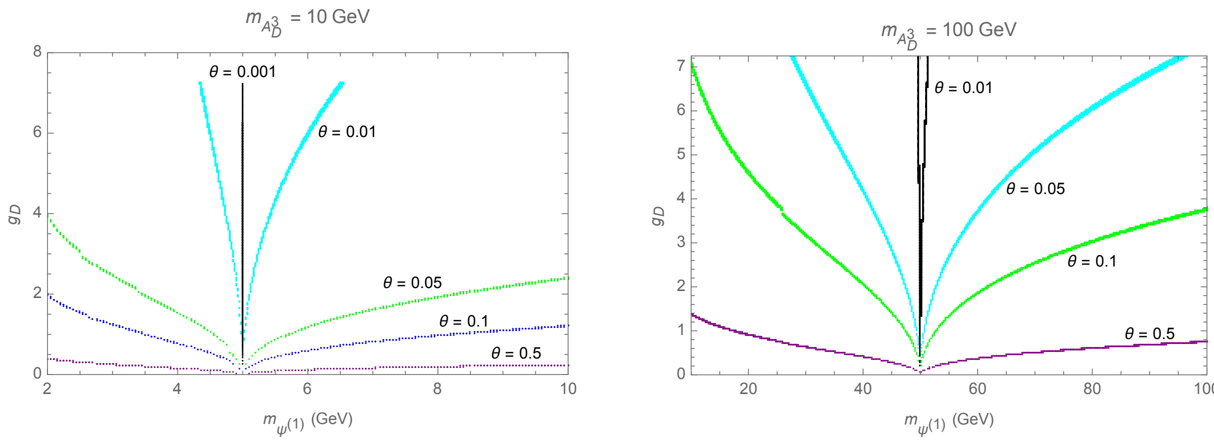

FIG. 2: Regions of the g_D-$m_{\psi^{(1)}}$ plane in which the dark matter relic density is within two standard deviations of the WMAP result 0.1186 ± 0.0020 [16], for fixed choices of $m_{A_D^3}$ and θ. The allowed bands are not perfectly smooth due to their dependence on g_*, which is not a continuous function. The point of minimum g_D corresponds to resonance annihilation, where $m_{\psi^{(1)}} = m_{A_D^3}/2$. Note that as θ decreases the range of $m_{\psi^{(1)}}$ in which g_D remains perturbative moves towards the resonance region.

that this is consistent with, for example, $v_{D1} = v_{D2} \approx 14$ GeV, $v_T \approx 49$ GeV, $\lambda_s \approx 0.85$, and $y_1 = y_2 \approx 0.06$, where the Yukawa couplings y_i were defined in Eq. (2.6). Similar statements can be made about other points on the allowed bands[3].

B. Direct Detection

The interactions that we have discussed to this point have involved leptons exclusively, but couplings to quarks that are generated at the loop level also have significant consequences.

[3] The scenario that we have considered assumes that communication between dark and visible sectors occurs primarily through the portal that we have proposed, involving mixing with vector-like leptons. It is of course possible to have scenarios in which communication is also significant through Higgs portal couplings or other mediators. The present model could therefore represent a subset of the parameters space of a more complicated model with other dark matter annihilation channels. There are also different parameter regions in the model as we have defined it where other annihilation channels become relevant, for example $\psi^{(1)}\overline{\psi^{(1)}} \to A_D^3 A_D^3$, when $m_{A_D^3} \leq m_{\psi^{(1)}}$. The results presented in this section demonstrate the effectiveness of the portal we have proposed in a region of parameter space where it provides the dominant contribution to the dark matter annihilation cross section due to mixing effects that would not be possible in models without flavor sequestering. This does not imply that other viable regions of parameter space are impossible.

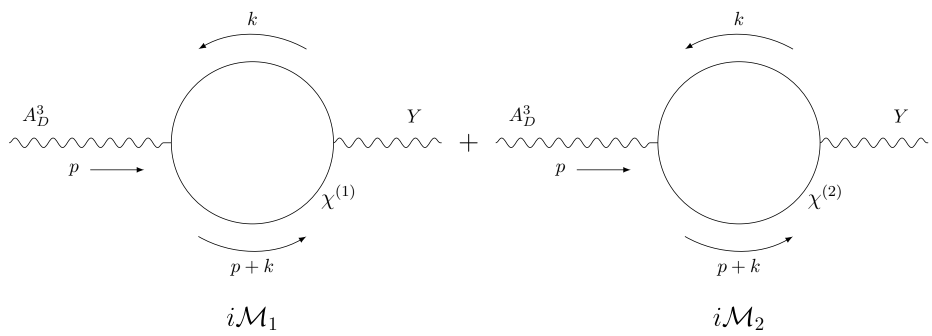

FIG. 3: Self energies leading to kinetic mixing between the third dark gauge boson $A_D^{(3)}$ and hypercharge Y after $\mathrm{SU}(2)_D$ is spontaneously broken.

In this section, we consider direct detection of the dark matter in the model via dark-matter-nucleon elastic scattering. The couplings to quarks arise after the $\mathrm{SU}(2)_D$ symmetry is spontaneously broken, since kinetic mixing between A_D^3 and hypercharge is then allowed, via an effective dimension-5 operator

$$\mathcal{L}_{eff} = X\,\mathrm{tr}\left(\langle H_T\rangle T^a A^a_{D\mu\nu}\right) Y^{\mu\nu} \quad , \tag{4.11}$$

where we have set the triplet Higgs to its vev, as per Eq. (2.4). Here, X is a constant with units of GeV^{-1} which is found by integrating out the "heavy" physics, *i.e.*, the χ fields, the only fields that are charged both under $\mathrm{SU}(2)_D$ and hypercharge $\mathrm{U}(1)_Y$. To proceed, we study the self-energy shown in Fig. 3, where $\chi^{(1)}$ and $\chi^{(2)}$ here represent the heavy mass eigenstates, whose mass eigenvalues are given approximately by $m_{\chi^{(1)}} = M_\chi - \delta$ and $m_{\chi^{(2)}} = M_\chi + \delta$ where $\delta \equiv \lambda_s'\, v_T/2$. (For the purposes of this estimate, we ignore mass mixing with the standard model lepton, which is a subleading correction.) The first diagram is given by

$$i\mathcal{M}_1 = -\frac{g_D g_Y}{2}\int \frac{d^4k}{(2\pi)^4}\frac{\mathrm{tr}\left[\gamma^\mu\left(\not{k} + m_{\chi^{(1)}}\right)\gamma^\nu\left(\not{k} + \not{p} + m_{\chi^{(1)}}\right)\right]}{[k^2 - m^2_{\chi^{(1)}} + i\epsilon][(k+p)^2 - m^2_{\chi^{(1)}} + i\epsilon]} \quad . \tag{4.12}$$

After carrying out this loop integral using dimensional regularization in $D = 4-\epsilon$ dimensions, the amplitude is

$$i\mathcal{M}_1 = -\frac{g_D g_Y}{8\pi^2}\int_0^1 dx\; x(1-x)\left(\frac{4}{\epsilon} - 2\gamma + 2\log(4\pi) - 2\log(\Delta_1)\right) i(g^{\mu\nu}p^2 - p^\mu p^\nu) \;\; , \tag{4.13}$$

where $\Delta_1 = m^2_{\chi^{(1)}} - x(1-x)p^2$. Since A_D^3 couples to the χ proportional to $\sigma^3/2$, the amplitude $i\mathcal{M}_2$ shown in Fig. 3 will differ from $i\mathcal{M}_1$ by a overall minus sign and the replacement of the $\chi^{(1)}$ by the $\chi^{(2)}$ mass. Hence, Δ_1 is replaced by $\Delta_2 = m^2_{\chi^{(2)}} - x(1-x)p^2$. Then, when these

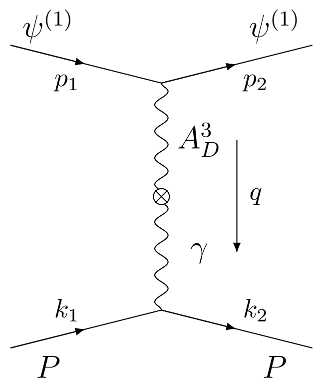

FIG. 4: The Feynman diagram for the scattering of the dark matter particles, $\psi^{(1)}$, off of protons, P, through kinetic mixing of the dark matter boson A_D^3 and the photon, γ.

two amplitudes are added together, all terms in the remaining Feynman parameter integral cancel, except for the terms that depend on the fermion masses:

$$i\mathcal{M}_1 + i\mathcal{M}_2 = i(g^{\mu\nu}p^2 - p^\mu p^\nu)\frac{g_D g_Y}{4\pi^2}\int_0^1 dx\; x(1-x)\log\left(\frac{\Delta_1}{\Delta_2}\right) \quad . \tag{4.14}$$

Assuming the mass splitting δ is small compared to the χ masses (which will turn out to be the case) the integrand can be expanded in δ. The leading order term can be found using $x(1-x)\log(\Delta_1/\Delta_2) \approx -\frac{4mx(1-x)}{m^2-x(1-x)p^2}\delta$. Moreover, we can also expand the result in powers of momentum, which can later be compared to a derivative expansion in the low-energy effective theory. We find

$$i\mathcal{M}_1 + i\mathcal{M}_2 = -i\frac{g_D g_Y \delta}{6\pi^2 M_\chi}(g^{\mu\nu}p^2 - p^\mu p^\nu) + \cdots \quad , \tag{4.15}$$

where the $\cdots$ represents terms involving higher powers of δ and p^2/M_χ^2. The result in Eq. (4.15) must be matched to a similar amplitude in the low-energy effective theory in which the χ fields have been integrated out. We identify this as the tree-level amplitude associated with the Eq. (4.11), treated as a two-point vertex,

$$i\mathcal{A} = iXv_T\left(p^2 g^{\mu\nu} - p^\mu p^\nu\right) \quad , \tag{4.16}$$

from which we conclude

$$X = -\frac{g_D g_Y \delta}{6\pi^2 M_\chi v_T} \quad . \tag{4.17}$$

Using Eqs. (4.11) and (4.17), we can now calculate the cross section for dark matter scattering off of nucleons. We will be working in the limit of low momentum transfer

$q \sim \mathcal{O}(100)$ MeV ($\ll M_\chi$), where the effective description is accurate and where scattering through the Z boson is suppressed by $q^2/m_Z^2 \sim 10^{-6}$ compared to the photon. Hence, we will consider kinetic mixing involving the photon only from here on. First, we consider the dark matter, $\psi^{(1)}$, scattering off of a quark, q_f, as in the diagram in Fig. 4, with the protons replaced by a quark of flavor f. This can be described by the effective dimension-six operator

$$\mathcal{L}_{eff,q} = C_f\, \overline{\psi^{(1)}}\gamma^\mu\psi^{(1)}\overline{q_f}\gamma_\mu q_f \quad . \tag{4.18}$$

In the full theory, this quark-dark matter scattering amplitude is

$$i\mathcal{M}_f = iXv_TQ_f\frac{g_D}{2}e\frac{1}{(q^2 - m^2_{A^3_D} + i\epsilon)(q^2 + i\epsilon)}\overline{\psi^{(1)}}\gamma^\mu\psi^{(1)}\overline{q_f}\gamma_\mu q_f \tag{4.19}$$

or, in the limit of $q^2 \ll m^2_{A_D^{(3)}}$,

$$i\mathcal{M}_f = -i\frac{Xv_TQ_f}{m^2_{A^3_D}}\frac{g_D}{2}e\overline{\psi^{(1)}}\gamma^\mu\psi^{(1)}\overline{q_f}\gamma_\mu q_f \quad . \tag{4.20}$$

From this, we conclude the coefficient C_f for quarks is

$$C_f = -\frac{g_d e X v_T Q_f}{2m^2_{A^3_D}} = \frac{g_D^2 e^2\delta Q_f}{12\pi^2 M_\chi m^2_{A^3_D}} \quad . \tag{4.21}$$

Of interest, however, is the effective interactions involving nucleons rather than quarks, which can be written

$$\mathcal{L}_{eff,N} = C_n\, \overline{\psi^{(1)}}\gamma^\mu\psi^{(1)}\overline{n}\gamma_\mu n + C_p\, \overline{\psi^{(1)}}\gamma^\mu\psi^{(1)}\overline{p}\gamma_\mu p \quad . \tag{4.22}$$

Using the fact that the quark vector currents are conserved, so that the spatial integral of the zeroth component is a quark number operator, one can match matrix elements of Eq. (4.18) between nucleon states with the same for Eq. (4.22), from which one concludes $C_n = C_u + 2C_d$ and $C_p = 2C_u + C_d$, for the neutron and proton, respectively. (There are no form factors as there would be for scalar quark operators.) Since the flavor dependence of the C_f comes only from the electric charge, the coefficient C_n and thus the scattering amplitude for $\psi^{(1)}$ off of neutrons are both zero. Therefore, the only relevant scattering is with the proton, for which

$$C_p = 2\,C_u + C_d = \frac{g_D^2 e^2\delta}{12\pi^2 M_\chi m^2_{A^3_D}} \quad . \tag{4.23}$$

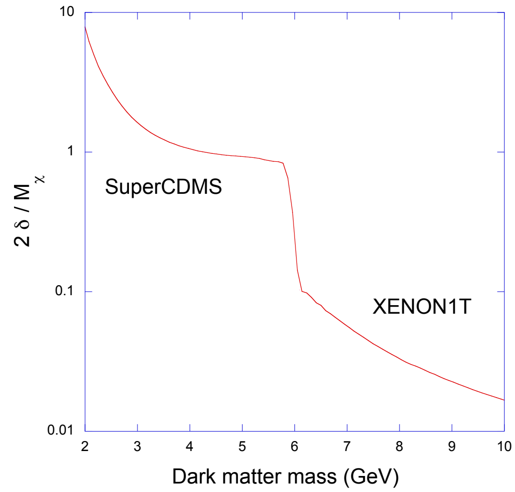

FIG. 5: Upper bound on the fractional mass splitting of the $\chi^{(1)}$ and $\chi^{(2)}$ fermions as a function of the mass of the dark matter particle, $\psi^{(1)}$, assuming $g_D = 0.3$ and $m_{A_D^3} = 10\,\text{GeV}$. The discontinuity in the curve reflects that the bounds on the dark-matter-nucleon elastic scattering cross section originate from the CDMSlite experiment [19] below $m_{\psi^{(1)}} \approx 6$ GeV, where the otherwise tighter bounds from the XENON1T experiment [18] do not exist.

Taking into account that the dark matter is non-relativistic and that momentum transfers are small, a straightforward calculation of the scattering cross section yields

$$\left\langle \sigma_{\psi^{(1)}p\to\psi^{(1)}p} \right\rangle = \frac{g_D^4 e^4 m_p^2 m_{\psi^{(1)}}^2}{576\pi^5 \left(m_p + m_{\psi^{(1)}}\right)^2 m_{A_D^3}^4} \left(\frac{2\delta}{M_\chi}\right)^2 \,, \tag{4.24}$$

where we have separated out the dependence on $2\delta/M_\chi$, the fractional mass spitting of the vector-like leptons. Since this splitting is a free parameter in our model, we can use the experimental bounds on the dark-matter-nucleon elastic scattering cross section to say something about the vector-like lepton spectrum.

Using experimental bounds on the cross section from XENON1T [18] and CDMSlite [19], we show bounds on the $\chi^{(1)}$-$\chi^{(2)}$ mass splitting for dark matter masses between 2 GeV and 10 GeV. The results of this calculation are shown in Fig. 5, where a dark coupling of $g_D = 0.3$ and a dark boson mass of $m_{A_D^3} = 10\,\text{GeV}$ have been used. For dark matter masses below approximately 6 GeV, the cross section bounds from CDMSlite are used, since no data from XENON1T is available in this region. Although there is CDMSlite data for dark matter masses above 6 GeV, these bounds are superceded by the stricter ones from XENON1T. For the range of $\psi^{(1)}$ masses in Fig. 5 that are affected by the XENON1T bounds, the masses of the charged fermions $\chi^{(1)}$ and $\chi^{(2)}$ are degenerate at the 1-10% level at minimum. This feature could be observed in collider searches for the vector-like leptons and possibly correlated with a dark matter direct detection signal.

V. CONCLUSIONS

We have presented a framework based on flavor symmetries that allows for a light portal sector of vector-like leptons connecting a dark sector to the standard model. To illustrate our approach, we considered an explicit, renormalizable non-Abelian dark SU(2) model which contains two vector-like fermion doublets. One of them, ψ, includes a dark matter candidate; the other doublet, χ, has the same electroweak quantum numbers as a right-handed electron, so that communication with the visible sector can occur via mass mixing. The ψ and χ fields communicate with each other via the dark gauge group, so that the dark matter may annihilate to standard model leptons. The dark SU(2) symmetry is spontaneously broken via a Higgs sector involving doublet and triplet fields. The doublet vacuum expectation value (vev) leads to mixing between the χ and standard model lepton fields, while the triplet vev splits the mass spectrum leaving a simple lower-energy theory consisting of the dark matter (the lightest ψ mass eigenstate) and the mediator (the third component of the SU(2) gauge multiplet). We identify a discrete flavor symmetry that allows mixing between the vector-like leptons χ and a single standard model lepton flavor exclusively; the remaining standard model lepton flavors may mix only with each other. This flavor sequestering eliminates lepton-flavor-violating effects, relaxing bounds on the vector-like lepton mass scale. As a consequence, mixing between the chosen lepton flavor and the χ can be large enough so that the correct relic density can be obtained exclusively via dark matter annihilation to lepton-anti-lepton pairs, for perturbative values of the dark gauge coupling. This is true even if no other significant annihilation channels are available.

The structure of our model avoids complications that would ensue if we tried to couple the dark gauge bosons directly to standard model fields, such as the necessity of including extraneous fermions to cancel chiral anomalies, or special Higgs representations to allow for acceptable standard model Yukawa couplings. Unlike some of the non-Abelian dark matter models appearing in the literature, the portal we present is renormalizable and completely specified, including the discrete flavor symmetries that control the pattern of mixing between exotic and standard model fermions states. The portal we define for communication between the dark and visible sectors can be lighter than in models without the flavor sequestering we have proposed and presents a well defined framework for answering phenomenological questions. In the present work, we showed that there are regions of the dark gauge coupling - dark matter mass plane where the correct relic density is obtained, and where current direct detection bounds are satisfied. The latter consideration also allowed us to conclude that the two heavy lepton mass eigenstates (roughly the two components of the χ doublet) are notably degenerate in mass (to keep kinetic mixing effects small), a feature that could be tested in collider searches for these states. This observation, together with the distinct lepton flavor structure of the χ decays, suggests that the collider signatures of the portal that we have proposed are worthy of future detailed investigation [14].

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