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Exotic Charges, Multicomponent Dark Matter and Light Sterile Neutrinos

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Generating small sterile neutrino masses via the same seesaw mechanism that suppresses active neutrino masses requires a specific structure in the neutral fermion mass matrix. We present a model where this structure is enforced by a new $U(1)'$ gauge symmetry, spontaneously broken at the TeV scale. In order not to spoil the neutrino structure, the additional fermions necessary for anomaly cancellations need to carry exotic charges, and turn out to form multicomponent cold dark matter. The active–sterile mixing then connects the new particles and the Standard Model—opening a new portal in addition to the usual Higgs- and kinetic-mixing portals—which leads to dark matter annihilation almost exclusively into neutrinos.

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I. INTRODUCTION

Neutrino oscillation experiments have by now firmly established the existence of neutrino oscillations and lepton flavor mixing, indicating that the Standard Model (SM) framework in particle physics has to be extended to include neutrino masses. Apart from the traditional three neutrino oscillation picture, the LSND [1] and MiniBooNE [2] short-baseline experiments suggest the presence of sterile neutrinos at the eV scale, which do not participate in the weak interaction but mix with active neutrinos with a mixing angle $\theta_s \sim \mathcal{O}(0.1)$ [3, 4]. Furthermore, recent re-evaluations of reactor anti-neutrino fluxes indicate that the previous reactor neutrino experiments had observed a flux deficit, which can in fact be interpreted by additional sterile neutrinos with masses at the eV scale (see Ref. [5] for an exhaustive overview of the field). Moreover, the light-element abundances from precision cosmology and Big Bang nucleosynthesis favor extra radiation in the Universe, which could be interpreted with the help of one additional sterile neutrino, albeit with a mass below eV [6].

From the theoretical point of view, the question then arises how the small sterile neutrino mass scale $\mathcal{O}$(eV) can be motivated compared to the electroweak scale of $\mathcal{O}$(100 GeV). Recall that the seesaw mechanism [7] is one of the most popular theoretical attempts to understand the smallness of active neutrino masses and also the baryon number asymmetry of the Universe. An obvious ansatz is then to use the same seesaw mechanism to suppress sterile neutrino masses. To this end, the right-handed neutrino content has to be extended compared to that in the simplest type-I seesaw mechanism, and a specific flavor structure, i.e., the minimal extended seesaw (MES), has to be employed [8, 9]. Explicitly, in the MES model, the SM fermion content is extended by adding three right-handed neutrinos ν_{Ri} (for $i = 1, 2, 3$) together with one singlet fermion S,‡ while the full Majorana mass matrix for the neutral fermions in the basis $(\nu_e, \nu_\mu, \nu_\tau, \nu^c_{R,1}, \nu^c_{R,2}, \nu^c_{R,3}, S^c)$ is given by

$$\mathcal{M}_{\mathrm{MES}} = \begin{pmatrix} 0 & m_D & 0 \\ m_D^T & M_R & m_S \\ 0 & m_S^T & 0 \end{pmatrix} . \tag{1}$$

After integrating out the heavy right-handed neutrinos $\nu_{R,i}$, one obtains three massive eigenstates with masses around M_R, three with masses around the eV scale and one massless neutrino. Taking a hierarchical spectrum $m_D < m_S \ll M_R$, one of the three massive eigenstates will be a sterile neutrino and its admixture with active neutrinos is suppressed by a factor $\mathcal{O}(m_D/m_S)$ (see Ref. [9] and Sec. III A for detailed discussions).

Note that the MES structure defined in Eq. (1) can be obtained with discrete flavor symmetries [9], under which the right-handed neutrinos and S carry different charges. In addition, the MES structure

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‡ To clear up potential confusion right away: all the $\nu_{R,j}$ and S_j introduced in this paper are just right-handed fermions, sometimes referred to as singlets. We denote the S_j with a different symbol than ν_{Ri} to emphasize that they are not the usual right-handed neutrinos from the seesaw mechanism, because they carry additional (hidden) quantum numbers and do therefore not partner-up with the active neutrinos.

could also be obtained in models with abelian symmetries. For example, one may introduce an extra $U(1)'$ symmetry, and take the three right-handed neutrinos $\nu_{R,i}$ to be neutral under $U(1)'$ (all SM particles are neutral too). One may then write down a bare Majorana mass matrix M_R for ν_R, which is unprotected by the electroweak or $U(1)'$ scale. The right-handed singlet S on the other hand carries a $U(1)'$ charge Y', and we further introduce an SM singlet scalar ϕ with charge $-Y'$. The gauge invariant coupling $\overline{S^c}\nu_R\phi$ then generates the m_S matrix in Eq. (1) after ϕ acquires a vacuum expectation value (VEV), while the Majorana mass for S (i.e., $\overline{S^c}S$) and a coupling to the active ν_L are still forbidden by the $U(1)'$ symmetry at the renormalizable level. Such a simple realization of MES however suffers from the problem of triangle anomalies, and can therefore only work as a global symmetry, whose spontaneous breaking would result in a massless Goldstone boson. While this might not be disastrous, more interesting phenomenology arises when the $U(1)'$ is promoted to a local symmetry. Consequently, one has to extend the model by additional chiral fermions so as to cancel the arising gauge anomalies. Along these lines, possible model constructions for sterile neutrinos in the $U(1)'$ framework have already been discussed in Refs. [10], using an effective-field-theory approach.

In this note we will work in the seesaw framework and discuss minimal renormalizable and anomaly-free $U(1)'$ models, which are spontaneously broken by just one additional scalar and reproduce the MES structure accounting for the $3+1$ or $3+2$ scheme of light sterile neutrinos. In particular, we will show that the additional singlet fermions employed for the anomaly cancellation turn out to be stable—due to remaining $\mathbb{Z}_N$ symmetries [11]—implying that they can be viewed as good candidates for the dark matter (DM) in the Universe. The remaining parts of this work are organized as follows: In Sec. II we explain the framework of exotic charges under a gauged $U(1)'$ and why they are needed to obtain the MES structure. In Sec. III we discuss in some detail the phenomenology of a specific example with one light sterile neutrino ($3+1$ scheme) and three stable DM candidates, with a focus on the novel effects inherent in our model. We briefly discuss other interesting examples of this framework in Sec. IV, including an extension to the $3+2$ case. Finally, we summarize our work in Sec. V.

II. DARK SYMMETRY

As already mentioned in the introduction, adding just one extra right-handed singlet S to the three ν_R's results in triangle anomalies if only S is charged under the extra $U(1)'$ symmetry. Instead of treating $U(1)'$ as a global symmetry, we gauge the $U(1)'$ in the rest of the work, and accordingly introduce additional singlet chiral fermions to cancel the anomalies. As we will see below, these new states need to decouple from the neutrino sector in order not to spoil the MES structure, and automatically lead to DM candidates without the need for additional stabilizing discrete symmetries.

For a gauged $U(1)'$ symmetry under which all SM particles are singlets, there are no mixed triangle anomalies, so anomaly freedom reduces to the equations

$$\sum_f Y'(f) = 0 \qquad \text{and} \qquad \sum_f (Y'(f))^3 = 0\,, \tag{2}$$

where f stands for our new right-handed fermions. In order to cancel the contribution from $S \equiv S_1$, more $U(1)'$ charged chiral fermions $S_{i\geq 2}$ have to be introduced. The solutions of Eq. (2) for $n = 2$ are simply given by $Y'(S_1) = -Y'(S_2)$. In this case, a bare mass term $m\overline{S_1^c}S_2$—unconstrained by any symmetry—can be constructed, which spoils the desired MES structure for light sterile neutrinos unless we make m small. There is no integer solution with $Y' \neq 0$ for $n = 3$ according to the famous Fermat theorem. In the case of $n = 4$, it is easy to prove that there is no phenomenologically interesting solution since two of the S_i must have $U(1)'$ charges of opposite sign and equal magnitude, inducing an unconstrained bare mass term as in the case of $n = 2$. For $n \geq 5$ however, there exist interesting non-trivial anomaly-free charge assignments—dubbed exotic charges hereafter—for example the set (10, 4, -9, 2, -7) for $n = 5$ [10, 12]. In order to make all new fermions massive at tree level with just one scalar ϕ, even more chiral singlets have to be introduced. For the $3+1$ scheme discussed in the main text, we add seven singlet fermions S_i to the model (the $3+2$ scheme discussed in Sec. IV A needs six). The charges of all the ten right-handed fermions are listed in Tab. I; they are of course by no means unique, but serve as a simple illustration of this framework. We further stress that at least three $U(1)'$ singlet right-handed neutrinos $\nu_{R,i}$ are needed in order to explain the observed light neutrino mass-squared differences Δm_{21}^2, Δm_{31}^2 and Δm_{41}^2, resulting in one massless active neutrino. This is however not a hard prediction of the MES scheme; adding a fourth ν_R (or even more) to the model makes all light neutrinos massive and does not qualitatively change or complicate the discussion below. Other interesting charge assignments with similar overall phenomenology are presented in Sec. IV.

	$\nu_{R,1}$	$\nu_{R,2}$	$\nu_{R,3}$	S_1	S_2	S_3	S_4	S_5	S_6	S_7	ϕ
Y'	0	0	0	11	−5	−6	1	−12	2	9	11

Table I: $U(1)'$ charge assignments of the right-handed fermions and the scalar ϕ.

In the scalar sector, we adopt only one SM singlet scalar ϕ with $U(1)'$ charge 11. We can then write down the following renormalizable couplings relevant for the neutrino masses

$$-\mathcal{L}_m = (m_D)_{ij}\overline{\nu_{L,i}}\nu_{R,j} + \frac{1}{2}(M_R)_{ij}\overline{\nu^c_{R,i}}\nu_{R,j} + w_i\phi^\dagger\,\overline{S^c_1}\nu_{R,i} + y_1\phi\,\overline{S^c_3}S_2 + y_2\phi\,\overline{S^c_4}S_5 + y_3\phi^\dagger\,\overline{S^c_6}S_7 + \text{h.c.}, \tag{3}$$

where appropriate sums over i and j are understood. The m_D terms stem from electroweak symmetry breaking using the usual SM Higgs doublet H, while w_i and y_i are Yukawa couplings. Absorbing phases into the S_j we can take y_j and one of the w_j to be real, while M_R can taken to be real and diagonal as well. Once ϕ acquires a VEV, the m_S matrix will be generated as $m_S = w_j\langle\phi\rangle$, leading to the MES structure in Eq. (1) (discussed in detail in Sec. III A). The other fermions $S_{2,3,4,5,6,7}$ decouple from ν_L, ν_R and S_1, and actually can be paired together to form three (stable) Dirac fermions $\Psi_{1,2,3}$, to be discussed in Sec. III C.

Let us briefly comment on a theoretical constraint on the model. An inherent problem in any gauge theory involving abelian factors is the occurrence of a Landau pole, i.e., a scale at which the gauge coupling becomes so large that our perturbative calculations break down. In our model, the one-loop beta function of the $U(1)'$ gauge coupling g' takes the form

$$\frac{\mathrm{d}}{\mathrm{d}\ln\mu}g' = \beta = \frac{{g'}^3}{16\pi^2}b = \frac{{g'}^3}{16\pi^2}\left[\frac{2}{3}\sum_j(Y'(S_j))^2 + \frac{1}{3}(Y'(\phi))^2\right], \tag{4}$$

so the Landau pole of g' appears around the scale

$$\Lambda_L \simeq \Lambda' \exp\left(\frac{8\pi^2}{b\,(g'(\Lambda'))^2}\right), \tag{5}$$

where Λ' characterizes the $U(1)'$ breaking scale. Inserting the $U(1)'$ charges given in Tab. I we find $b = 315$, whereas for the $3+2$ scheme from Tab. II (which will be discussed later on in Sec. IV A) we have $b = 75$. For $\Lambda' \simeq 1$ TeV and $\Lambda_L \gtrsim M_{\rm Pl} \simeq 10^{19}$ GeV, one obtains the constraints $g'(\Lambda') \lesssim 0.08$ for the $3+1$ case and $g'(\Lambda') \lesssim 0.17$ for $3+2$ case. Alternatively, if we take the cutoff scale of the model to be the right-handed neutrino mass scale, i.e., $\Lambda_L \gtrsim M_R \sim 10^{14}$ GeV, these bounds relax to $g'(\Lambda') \lesssim 0.1$ and $g'(\Lambda') \lesssim 0.2$ for $3+1$ and $3+2$, respectively. These upper bounds are stricter than the naive perturbativity bound ${g'}^2/4\pi \lesssim \mathcal{O}(1)/\max(Y')^2$.

III. PHENOMENOLOGICAL CONSEQUENCES

In this section we discuss the phenomenology of our new particles, with an emphasis on the novel effects in our framework.

A. Neutrino Masses and Active–Sterile Mixing

Let us consider the neutrino masses and active–sterile mixing in the $U(1)'$ model. After the breaking of the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)'$ to $SU(3)_C \times U(1)_{\rm EM}$, the full 13×13 mass matrix for the neutral fermions in the basis $\nu = (\nu_{L,1}, \nu_{L,2}, \nu_{L,3}, \nu^c_{R,1}, \nu^c_{R,2}, \nu^c_{R,3}, S^c_1, S^c_2, S^c_3, S^c_4, S^c_5, S^c_6, S^c_7)$ reads

$$\mathcal{M} = \begin{pmatrix} (\mathcal{M}_{\rm MES})_{7\times 7} & 0 \\ 0 & (\mathcal{M}_S)_{6\times 6} \end{pmatrix}, \tag{6}$$

where the matrix $\mathcal{M}_{\rm MES}$ reproduces the MES structure derived in Ref. [8, 9] (although our definition for m_S differs by transposition), while $\mathcal{M}_S$ denotes the mass matrix of S_{2-7}, explicitly given as

$$\mathcal{M}_S = \begin{pmatrix} 0 & y_1\langle\phi\rangle & 0 & 0 & 0 & 0 \\ y_1\langle\phi\rangle & 0 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & y_2\langle\phi\rangle & 0 & 0 \\ 0 & 0 & y_2\langle\phi\rangle & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & y_3\langle\phi\rangle \\ 0 & 0 & 0 & 0 & y_3\langle\phi\rangle & 0 \end{pmatrix} . \tag{7}$$

Obviously S_{2-7} decouple from the neutrino sector and can no longer be interpreted as right-handed neutrinos, because they do not partner-up with the SM neutrinos. The bare mass term M_R is unrestricted and can be large, as in the canonical seesaw case. We will consider this possibility here by setting $M_R \gg m_D, m_S$, which leads to the effective low-energy neutrino mass matrix

$$\mathcal{M}_\nu^{4\times4} \simeq -\begin{pmatrix} m_D M_R^{-1} m_D^T & m_D M_R^{-1} m_S \\ m_S^T M_R^{-1} m_D^T & m_S^T M_R^{-1} m_S \end{pmatrix} , \tag{8}$$

for (ν_L, S_1^c).[§] Such a mass matrix can be diagonalized by means of a unitary transformation as $\mathcal{M}_\nu^{4\times4} = V\,{\rm diag}(m_1, m_2, m_3, m_4)\,V^T$. Phenomenologically, the most interesting situation arises for $m_S \gg m_D$, since the hierarchical structure of $\mathcal{M}_\nu^{4\times4}$ allows us to apply the seesaw expansion once more, and arrive at the sterile neutrino mass

$$m_4 \simeq -m_S^T M_R^{-1} m_S \, , \tag{9}$$

together with the mass matrix for the three active neutrinos

$$\mathcal{M}_\nu^{3\times3} \simeq -m_D M_R^{-1} m_D^T + m_D M_R^{-1} m_S \,(m_S^T M_R^{-1} m_S)^{-1} m_S^T M_R^{-1} m_D^T = U\,{\rm diag}(m_1, m_2, m_3)\,U^T \, , \tag{10}$$

diagonalized by U. The 4×4 unitary mixing matrix V is approximately given by

$$V \simeq \begin{pmatrix} (1-\frac{1}{2}RR^\dagger)U & R \\ -R^\dagger U & 1-\frac{1}{2}R^\dagger R \end{pmatrix} \tag{11}$$

with the active–sterile mixing vector $R = m_D M_R^{-1} m_S\,(m_S^T M_R^{-1} m_S)^{-1} = \mathcal{O}(m_D/m_S)$. As a rough numerical estimation, for $m_D \sim 10^2$ GeV, $m_S \sim 5\times10^2$ GeV and $M_R \sim 2\times10^{14}$ GeV, one obtains the active neutrino mass scale $m_\nu \sim 0.05$ eV, the sterile neutrino mass scale $m_s \sim 1.3$ eV together with $R \simeq 0.2$. This is in good agreement with the current global-fit data, i.e., $|R_1| \simeq 0.15$ and $\Delta m_{41}^2 \simeq 1.8\ {\rm eV}^2$ for the $3+1$ scheme [4].

For Yukawa couplings of order one, the observed large active–sterile mixing implies the scaling $\langle\phi\rangle/\langle H\rangle \sim 5$–$10$. The new physics scale around TeV is hence not tuned to make LHC phenomenology most interesting, but comes directly from the neutrino sector. Actually—even though we obtain the magic TeV scale—the LHC implications of our model are rather boring, as we only expect small mixing effects in the Higgs and Z-boson interactions, to be discussed in the next section.

Let us briefly comment on thermal leptogenesis [13] in our framework. In principle, the additional singlet fermions may spoil the ordinary picture of leptogenesis since the right-handed neutrinos might predominately decay to sterile neutrinos instead of active neutrinos. This drawback can be easily circumvented here by choosing the coupling of the lightest right-handed neutrino $\nu_{R,1}$ to the new states to be small, i.e., $w_1 \ll w_{2,3}$. This will not modify the desired MES structure in the neutrino sector, but sufficiently increase the branching ratio of $\nu_{R,1}$ into SM particles, so standard thermal leptogenesis ensues.

B. Bosonic Sector

In this subsection we will briefly summarize the behavior of ϕ and Z'. The scalar potential W in our model takes a simple form

$$W = -\mu_H^2|H|^2 + \lambda_H|H|^4 - \mu_\phi^2|\phi|^2 + \lambda_\phi|\phi|^4 + \delta\,|H|^2|\phi|^2 \, , \tag{12}$$

[§] On a more fundamental level, one can integrate out the heavy right-handed neutrinos ν_R at energies $E \ll M_R$ to generate the effective dimension-five Weinberg operators $(m_D)_{ij}(m_D)_{kj}\overline{L}_i\tilde{H}H^\dagger\tilde{L}_k/(\langle H\rangle^2\,(M_R)_{jj})$, $w_i^2\,\phi^2\overline{S}_1S_1^c/(M_R)_{ii}$ and $(m_D)_{ij}w_j\,\overline{L}_i\tilde{H}S_1\phi^\dagger/(\langle H\rangle\,(M_R)_{jj})$, which were the starting point in Refs. [10].

where H denotes the SM Higgs doublet, and ϕ can be decomposed as $\phi = (\mathrm{Re}\,\phi + \mathrm{i}\,\mathrm{Im}\,\phi)/\sqrt{2} = (\langle \mathrm{Re}\,\phi\rangle + \varphi + \mathrm{i}\,\mathrm{Im}\,\phi)/\sqrt{2}$. After symmetry breaking and in unitary gauge, $\mathrm{Im}\,\phi$ is absorbed by the Z' boson, giving it a mass $M_{Z'} = |11 g' \langle\phi\rangle|$. Due to the δ term in the scalar potential, we have a generic mixing between the remaining real field φ and the neutral scalar h contained in H, i.e.,

$$\begin{pmatrix} h_1 \\ h_2 \end{pmatrix} = \begin{pmatrix} \cos\theta & -\sin\theta \\ \sin\theta & \cos\theta \end{pmatrix} \begin{pmatrix} h \\ \varphi \end{pmatrix}, \tag{13}$$

where h_1 and h_2 are the physical mass eigenstates, and the mixing angle θ is given by

$$\sin 2\theta = \frac{\delta \langle\phi\rangle\langle H\rangle}{\sqrt{\left(\lambda_\phi \langle\phi\rangle^2 - \lambda_H \langle H\rangle^2\right)^2 + \left(\delta \langle H\rangle\langle\phi\rangle\right)^2}}\,. \tag{14}$$

A non-zero δ—and hence θ—opens the well-known Higgs portal [14] for the DM production/annihilation, which will be discussed in the next section.

The Higgs portal $|\phi|^2 |H|^2$ aside, there is one more renormalizable, gauge-invariant operator that will induce a coupling between the SM and DM sectors, namely the kinetic-mixing operator $\sin\xi\, F_Y^{\mu\nu} F'_{\mu\nu}$ [15]. This off-diagonal kinetic term involving the hypercharge and $U(1)'$ field strength tensors will induce a coupling of the physical Z' boson to the hypercharge current. The relevant phenomenology of the resulting interaction between the SM and DM particles can be found in e.g. Refs. [16, 17].

C. Dark Matter

As we mentioned before, the singlet fermions $S_{2,3,4,5,6,7}$ are the DM candidates in our model. To see this point, we write down the full Lagrangian for the right-handed singlets S_2 and S_3:

$$\mathcal{L}_{S_{2,3}} = \mathrm{i}\overline{S}_2 \gamma^\mu (\partial_\mu - \mathrm{i}(-5g') Z'_\mu) S_2 + \mathrm{i}\overline{S}_3 \gamma^\mu (\partial_\mu - \mathrm{i}(-6g') Z'_\mu) S_3 + y_1 (\phi\, \overline{S}_3^c S_2 + \mathrm{h.c.})\,. \tag{15}$$

By defining the Dirac field $\Psi_1 \equiv S_2 + S_3^c$, the above Lagrangian can be rewritten in unitary gauge as

$$\mathcal{L}_{S_{2,3}} = \mathrm{i}\overline{\Psi}_1 \gamma^\mu \partial_\mu \Psi_1 + g' Z'_\mu \overline{\Psi}_1 \gamma^\mu \left(\frac{(-5)-(-6)}{2} + \frac{(-5)+(-6)}{2}\gamma_5 \right) \Psi_1 + \frac{y_1}{\sqrt{2}} \left(\langle \mathrm{Re}\,(\phi)\rangle + \varphi \right)\, \overline{\Psi}_1 \Psi_1\,. \tag{16}$$

After spontaneous symmetry breaking, Ψ_1 acquires a Dirac mass $M_1 \equiv -y_1 \langle \mathrm{Re}\,\phi\rangle/\sqrt{2}$. Similarly, we can define $\Psi_2 \equiv S_4 + S_5^c$ and $\Psi_3 \equiv S_6 + S_7^c$ for $S_{4,5}$ and $S_{6,7}$, and obtain altogether the DM Lagrangian

$$\begin{aligned} \mathcal{L}_{\mathrm{DM}} = \sum_{j=1,2,3} \Bigg[& \mathrm{i}\overline{\Psi}_j \gamma^\mu \partial_\mu \Psi_j - M_j \overline{\Psi}_j \Psi_j - \frac{M_j}{\langle \mathrm{Re}\,(\phi)\rangle}\, \varphi\, \overline{\Psi}_j \Psi_j \\ & + \frac{g'}{2} Z'_\mu \overline{\Psi}_j \gamma^\mu \left[(Y'_{2j} - Y'_{2j+1}) + (Y'_{2j} + Y'_{2j+1})\gamma_5 \right] \Psi_j \Bigg]\,, \end{aligned} \tag{17}$$

where we defined $Y'_j \equiv Y'(S_j)$. The stability of these fields will be discussed below, but let us first take a look at the interactions involving the will-be sterile neutrino S_1, given by the Lagrangian

$$\mathcal{L}_{S_1} = \mathrm{i}\overline{S}_1 \gamma^\mu (\partial_\mu - \mathrm{i}(11g') Z'_\mu) S_1 + \left(w_i \phi^\dagger\; \overline{S_1^c} \nu_{R,i} + \mathrm{h.c.} \right). \tag{18}$$

The important part is the Z' interaction, as it allows for the annihilation $\Psi_i \Psi_i \to Z' \to S_1 S_1$. Since the physical sterile neutrino $\nu_s \equiv \nu_4$ consists mainly of S_1, but contains a not-too-small part of the active neutrinos $\nu_{e,\mu,\tau}$, this process connects the DM to the SM sector. Specifically, this "neutrino portal" takes the form

$$\begin{aligned} \mathcal{L}_{\nu-\mathrm{portal}} = \frac{g'}{2} Z'_\mu \Bigg[& \overline{\Psi}_1 \gamma^\mu (1 - 11\gamma_5) \Psi_1 + \overline{\Psi}_2 \gamma^\mu (13 - 11\gamma_5) \Psi_2 + \overline{\Psi}_3 \gamma^\mu (-7 + 11\gamma_5) \Psi_3 \\ & + 11 \sum_{i,j=1}^{4} V_{4i}^* V_{4j}\, \left(\overline{\nu}_i \gamma^\mu \gamma_5 \nu_j + \overline{\nu}_i \gamma^\mu \nu_j \right) \Bigg]\,, \end{aligned} \tag{19}$$

where the four light mass eigenstates ν_j are written as Majorana spinors and the unitary matrix V is defined in Eq. (11).

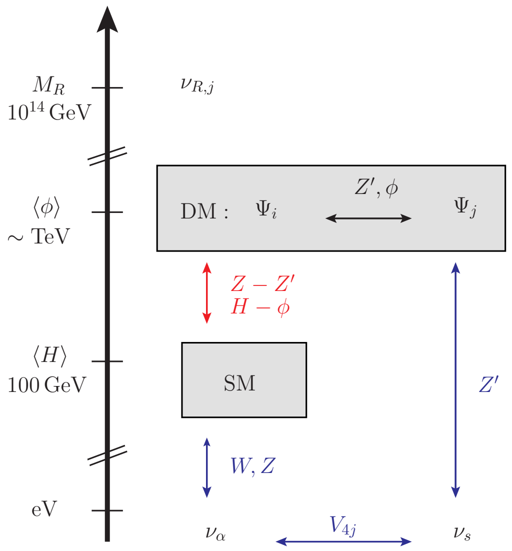

Figure 1: Visualization of the different scales in our framework, as well as the relevant interaction channels. The red connection between the DM and SM sectors represents the well-known kinetic-mixing and Higgs portals (based on vector and scalar mixing respectively), while the blue interactions are relevant for the neutrino portal (based on fermion mixing). Interactions with the $\nu_{R,j}$ are highly suppressed and not shown.

We further note that the heavier DM particles can also convert to the lighter ones, i.e., $\Psi_i\Psi_i \to \Psi_j\Psi_j$ via the s-channel exchange of Z' or ϕ. Moreover, Ψ_i may also annihilate to Z' and ϕ, which can enhance the total annihilation cross section significantly.

The model content and relevant scales are illustrated in Fig. 1. The (self-interacting) DM sector couples to the SM just like all models with a dark symmetry $U(1)_{\rm DM}$, namely through scalar mixing (Higgs portal, parameterized through δ) and vector mixing (kinetic-mixing portal, parameterized through ξ). However, due to the gauge interactions of the DM with the sterile neutrinos, a new portal through fermion mixing (neutrino portal) opens up in our model. Since this portal is not often discussed in the literature (see however Refs. [18, 19]), we will focus on it in the remainder of this paper.

1. *Stability*

It is obvious that the Ψ_i fields in Eq. (17) are stable, since there exists an accidental global $U(1)^3$ symmetry shifting the phases of Ψ_i. The occurrence of several stable DM particles (multicomponent DM) results in numerous interesting effects (see Refs. [20] for some early work). The underlying reason for the stability in our case is the remaining exact $\mathbb{Z}_{11}$ symmetry after the spontaneous breakdown of the $U(1)'$. The Ψ_j form representations under this discrete gauge group with charges 6, 1 and 2 (modulo 11), which stabilizes at least the lightest of them, even when higher-dimensional operators are considered.

While our model is renormalizable, we expect it to be only valid up to a certain cutoff scale Λ, either because quantum gravity takes over, or because sooner or later we will hit the $U(1)'$ Landau pole—as discussed in Sec. II. At the cutoff scale, higher-dimensional operators might be generated, and in our models these will always include $\phi^2\overline{S}_1 S_1^c/\Lambda$ and the Weinberg operator [21] for ν_L-Majorana masses. Taking $\Lambda \sim M_{\rm Pl}$ does not destroy the discussed MES structure if $\langle\phi\rangle \lesssim 10\,$TeV. For the charge assignment here, there are also dimension-six operators like $\overline{S}_2^c S_4 \overline{S}_6 S_3^c/\Lambda^2$, which break the global $U(1)^3$ to a $U(1)$ symmetry, so only one stable Dirac fermion survives. However, since these operators are highly suppressed for $\Lambda \sim M_{\rm Pl}$, the resulting lifetimes are typically longer than the age of the universe, and thus we will not include them in our discussions below, but take all three Ψ_j to be independently stable.

2. *Relic Density and Thermal History*

We will now discuss the interplay of the three portals (Higgs-, kinetic-mixing-, and neutrino portal) and identify some valid regions in the parameter space where the correct relic density for Ψ_j can be obtained. Note that the mixing parameters δ and ξ are the only new physics parameters we assume to be small in this paper, all other couplings are somewhat "natural". We restrict ourselves to small mixing parameters solely for simplicity, as larger values lead to very constrained effects, see Refs. [14, 16, 17, 22]. A more detailed parameter scan (including thermal leptogenesis) will be performed in a separate publication [23].

We only consider freeze-out scenarios. Also note that we always end up with a thermalized sterile neutrino at the epoch of neutrino decoupling, so the usual cosmological bounds on $N_{\rm eff}$ and $\sum m_\nu$ hold [5]. This is to be expected in models with light sterile neutrinos, and can be solved on the astrophysics side, or by choosing a smaller-than-eV mass for the sterile neutrino.

Case A: $\delta, \xi = 0$. To check the validity of the neutrino portal, we first turn off the Higgs- and kinetic-mixing portals by setting $\delta = \xi = 0$ (or at least small enough to be negligible). In this case, the only connection between the new physics sector and the SM comes from active–sterile mixing, or, at a more fundamental level, from the exchange of heavy right-handed neutrinos. Integrating out the ν_R gives for example the operator $LHS_1\phi/M_R$ (using order one Yukawa couplings), which gives a rough scattering rate for $LH \leftrightarrow S_1\phi$ around $\sim T^3/M_R^2$—to be compared to the expansion rate in the early universe $\sim \sqrt{g_*}T^2/M_{\rm Pl}$—which puts all particles in equilibrium above $T \gtrsim 10^{10}\,$GeV.¶ Below that temperature, the two sectors SM and DM (consisting of Z', ϕ and S_j) evolve independently, while the temperature decreases due to expansion of the Universe in both sectors. Nothing really happens until $T \sim$ TeV, when the Ψ_j freeze-out occurs. To deplete the Ψ_j abundance fast enough, we can make use of the neutrino portal, i.e., the annihilation of the lightest Ψ_j into ν_s around the Z' resonance (assuming that the heavier Ψ_j annihilate sufficiently fast into the lightest Ψ_j, see Fig. 2 for illustrations).

After freeze-out, we have overall three decoupled sectors—SM, Ψ_j and ν_s—all with different temperatures. Above (active) neutrino decoupling, the Universe was radiation dominated, so only the temperature of ν_s and the relativistic degrees of freedom in the SM sector are of interest and will be calculated now. Using conservation of entropy in the two sectors SM and DM, we have the equalities

$$g_*^{\rm SM} T^3 a^3 \Big|_{t_{\rm sep}} = g_*^{\rm SM} T_{\rm SM}^3 a^3 \Big|_{t_{\rm f}} \qquad \text{and} \qquad g_*^{\rm DM} T^3 a^3 \Big|_{t_{\rm sep}} = g_*^{\rm DM} T_{\rm DM}^3 a^3 \Big|_{t_{\rm f}} , \tag{20}$$

where g_*^X denotes the effective number of relativistic degrees of freedom in sector X, a the scale factor, $t_{\rm sep}$ the time when the two sectors just separated from equilibrium (i.e., at temperatures around $10^{10}\,$GeV) and t_f the final time we are interested in, namely close to active neutrino decoupling (e.g. when $T_{\rm SM} \sim 10\,$MeV). At t_f, the SM sector consists of photons, electrons and neutrinos, while the DM sector only has the relativistic $S_1 \sim \nu_s$, so we find

$$T_{\nu_s}/T_{\rm SM}\Big|_{t_{\rm f}} = \left(\frac{g_*^{\rm DM}(t_{\rm sep})}{g_*^{\rm DM}(t_f)}\frac{g_*^{\rm SM}(t_f)}{g_*^{\rm SM}(t_{\rm sep})}\right)^{1/3} = \left(\frac{65/4}{7/4}\frac{43/4}{427/4}\right)^{1/3} \simeq 0.98\,. \tag{21}$$

Ignoring active–sterile oscillations, this would make the sterile neutrinos slightly colder than the active ones at decoupling, alleviating cosmological constraints to some degree (the one sterile neutrino effectively contributes only $\Delta N_{\rm eff} = (T_{\nu_s}/T_{\rm SM})^4 \simeq 0.92$ additional neutrinos to the energy density). However, for the sterile neutrino parameters relevant for the short-baseline anomalies, i.e. $m_s \sim$ eV, $\theta_s \sim 0.1$, active–sterile oscillations will become effective around $T \sim 100\,$MeV–1 MeV [24], once again connecting the SM bath and ν_s and thus thermalizing the sterile neutrino at neutrino decoupling. Note that the usual discussions of active–sterile oscillations at these temperatures are not readily applicable, as our model starts with abundant ν_s and self-interactions mediated by Z' (freezing out around $T_{\rm DM} \sim 10\,$MeV). In any case, the cosmological bound on relativistic degrees of freedom is expected to be approximately valid in our model and will be discussed in more detail in a separate paper [23].

Case B: $\delta \neq 0$. Let us open the Higgs portal. The thermal evolution is similar to case A, but values $\delta \gtrsim 10^{-7}$ will put ϕ in equilibrium with the SM at temperatures below $T \sim 10\,$TeV, because the scattering rate $hh \leftrightarrow \phi\phi$ goes with $\delta^2 T/4\pi$ [17]. ϕ and the rest of the DM sector (Z', Ψ_j and ν_s) are in equilibrium through $U(1)'$ gauge interactions (for not too small gauge coupling g'), so SM and DM are in equilibrium around DM freeze-out. For the freeze-out we can again use the neutrino portal, i.e., resonant annihilation $\Psi\Psi \to Z' \to \nu_s\nu_s$. As ϕ and Z' go out of equilibrium around the same time, the connection between the SM sector and ν_s is severed and the two evolve independently for a while, until they are reconnected around $T \sim 10\,$MeV by active–sterile neutrino oscillations.

In a different region of parameterspace, we can make use of the resonant annihilation of DM into SM particles via scalars, i.e., the Higgs portal in the way it is intended. The discussion is then completely analogous to other $U(1)_{\rm DM}$ models, so we refer the interested reader to Ref. [22] for a recent evaluation.

¶ We assume a sufficiently high reheating temperature after inflation.

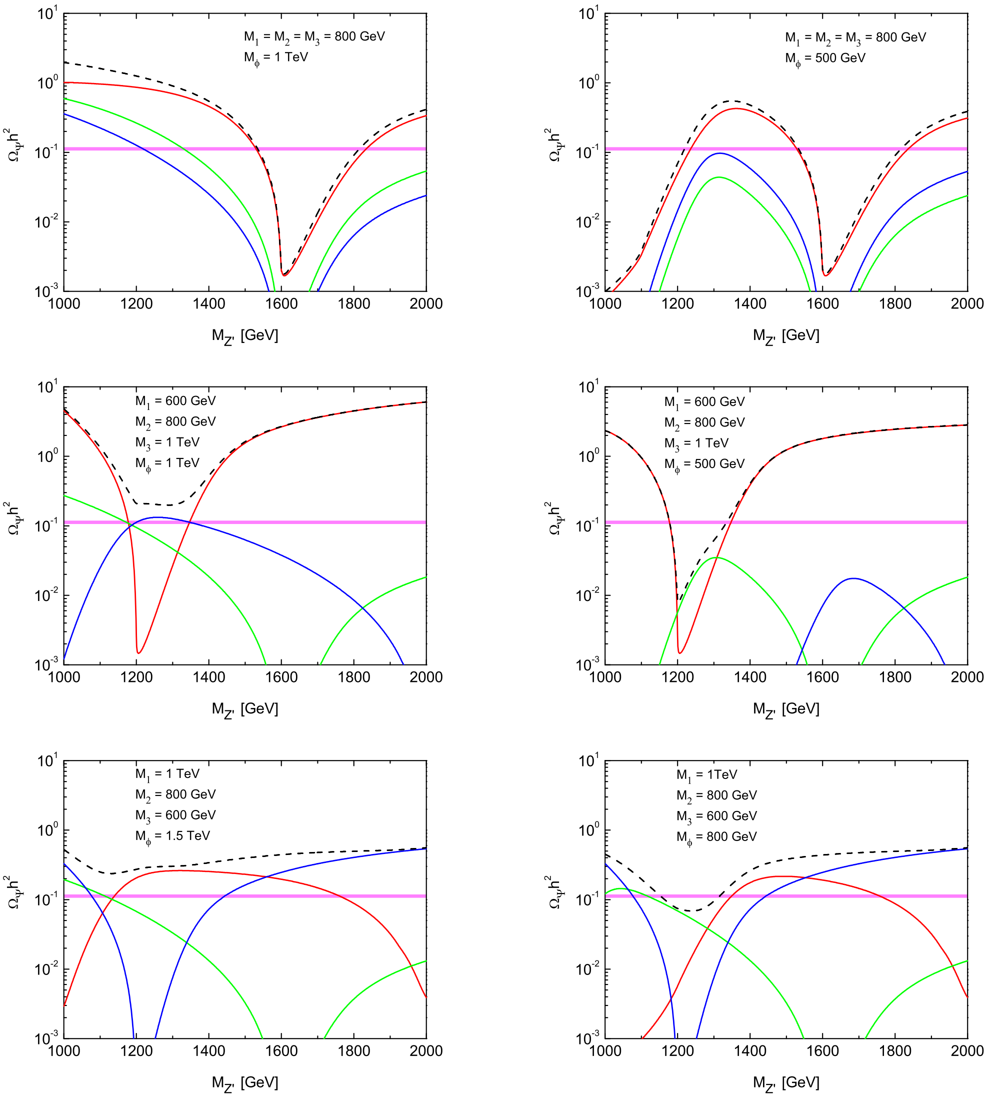

Figure 2: $\Omega_\Psi h^2$ versus the Z' mass $M_{Z'}$ for degenerate (top panels) and hierarchical (middle and bottom panels) DM masses. The VEV was set to $\langle\phi\rangle = 1.5\,\text{TeV}$ and the scalar mass is labeled on the plot. The red, green and blue lines show the relic density of Ψ_1, Ψ_2 and Ψ_3 respectively, while the black dashed line gives the full $\Omega_\Psi h^2 \equiv \sum_j \Omega_{\Psi_j} h^2$. The horizontal pink band represents the relic density measured by WMAP [25] (1σ range).

Case C: $\xi \neq 0$. A very similar discussion can be made for an open kinetic-mixing portal. Again small values $\xi \gtrsim 10^{-7}$ suffice to reach thermal equilibrium of the SM and DM sectors, e.g. through scattering $Zh \leftrightarrow Z'h$. The thermal evolution then closely resembles that of case B, with some minor differences: The Z'–Z mixing couples ν_s to the SM, so Z' interactions keep ν_s thermalized a while longer before it decouples and finally reconnects with the SM. Furthermore, the DM annihilation around the Z' resonance contains a small branching ratio into SM particles.

In order to illustrate the feasibility of the DM candidates via the neutrino portal, we implement the model in micrOMEGAs [26] and evaluate the relic density of DM particles Ψ_i. Here we assume δ and ξ to be large enough to thermalize the SM and DM sectors at DM freeze-out, but small enough to be negligible in the numerical calculation of the relic density, i.e., we only consider DM freeze-out using the neutrino portal. The scalar VEV is taken to be $\langle\phi\rangle = 1.5$ TeV as an example. The gauge coupling g' are therefore obtained from the relation of Z' mass and $\langle\phi\rangle$. As shown in the upper panels of Fig. 2, a resonance appears at $M_\Psi \simeq M_{Z'}/2$, and the relic density $\Omega_\Psi h^2 \simeq 0.1$ measured by WMAP [25] can be achieved. In the degenerate case (i.e., $M_1 \simeq M_2 \simeq M_3$), the Ψ_1 contribution to $\Omega_\Psi h^2$ is dominating because it has the smallest Z' coupling. Moreover, in case of a small scalar mass, e.g., $M_\phi = 500$ GeV, a new channel $\Psi\Psi \to Z'\phi$ is open for light Z', which is observed from the upper-right panel of Fig. 2. For the case of non-degenerate spectrum (i.e., $M_1 \neq M_2 \neq M_3$), the most significant contribution to the relic

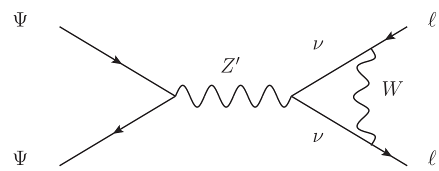

Figure 3: Coupling of DM to leptons.

density may come from either Ψ_1, Ψ_2 or Ψ_3, depending on the specific fermion spectrum as well as the scalar and vector masses. As can be seen in the middle and lower panels of Fig. 2, the Ψ_1 contribution to the relic density typically dominates, but there exist model parameters that make Ψ_2 or Ψ_3 the main DM particle.

3. Direct and Indirect Detection

The neutrino portal discussed so far does not lead to any direct detection signals, because the cross sections are highly suppressed. Loop processes connecting Ψ to SM fermions, e.g. as in Fig. 3, vanish in case of degenerate active–sterile masses, so these amplitudes are suppressed by tiny factors like $\Delta m_{41}^2/\mathcal{O}(100\,\mathrm{GeV})^2 \sim 10^{-22}$.

Indirect detection might naively be more fruitfull, because the annihilation of the Ψ_j in the Galactic Center or halo leads to two back-to-back neutrinos with energies $\simeq M_j$ (whichever Ψ_j is sufficiently abundant), which is an ideal signal for neutrino telescopes like IceCube.** However, since we considered Ψ_j to be a thermal relic, the self-annihilation cross section is set by the relic density, which is too small to be probed [27]—even though the branching ratio into neutrinos is $\simeq 100\%$, so the signal is as clear as it gets.

Direct and indirect detection measurements are of course sensitive to the Higgs- and kinetic-mixing portal parameters δ and ξ, as discussed in the literature (e.g. Ref. [19] discusses the Higgs portal in a framework similar to the neutrino portal).

IV. OTHER INTERESTING CHARGE ASSIGNMENTS

Having focussed on one specific example using the charges from Tab. I, we will now briefly present other charge assignments with interesting phenomenology. In all cases, we only introduce one additional scalar ϕ, so the results concerning scalar and vector interactions remain unchanged—different numerical values for the charges aside. Only the sterile neutrino and dark matter sector will be slightly modified.

A. More Light Sterile Neutrinos

The introduction of $n \geq 2$ light sterile neutrinos ($3+n$ scheme) increases the number of new parameters and most importantly allows for CP-violation in the effective oscillation analysis [28]. This feature can significantly improve the fit to neutrino oscillation data and has been studied extensively [3, 4, 29]. Note that the tension with standard ΛCDM cosmology typically worsens, depending on the used data sets [30].

We can easily modify the above $U(1)'$ framework to accommodate the $3+2$ MES scheme, by choosing different charges for the ten right-handed neutrinos: we need at least one more neutral $\nu_{R,4}$ to generate the necessary light mass squared differences. Now we have to find charges that treat two of S_i the same (without loss of generality S_1 and S_2), i.e., $Y'(S_1) = Y'(S_2)$, so these will become our two light neutrinos after coupling them to a scalar ϕ. We can once again find exotic charges in such a way that the decoupled S_j become massive by coupling to the same scalar, the magic number for this to happen seems to be six. See Tab. II for a valid anomaly-free charge assignment with the desired properties (also used in Ref. [10]). After breaking the $U(1)'$ and the electroweak symmetry, the 13×13 mass matrix for the neutral fermions

** The DM-nucleon cross section in our model is too small to efficiently capture DM inside the Sun or Earth, so we have to rely on astrophysical objects with high DM density.

	$\nu_{R,1}$	$\nu_{R,2}$	$\nu_{R,3}$	$\nu_{R,4}$	S_1	S_2	S_3	S_4	S_5	S_6	ϕ
Y'	0	0	0	0	-5	-5	-1	6	2	3	5

Table II: Exotic $U(1)'$ charge assignments of the right-handed fermions and the scalar ϕ to obtain the $3+2$ MES scheme.

takes the form

$$\mathcal{M} = \begin{pmatrix} (\mathcal{M}_{\rm MES})_{9\times 9} & 0 \\ 0 & (\mathcal{M}_S)_{4\times 4} \end{pmatrix} , \tag{22}$$

where the 9×9 matrix $\mathcal{M}_{\rm MES}$ in the basis $(\nu_e, \nu_\mu, \nu_\tau, \nu_{R,1}^c, \nu_{R,2}^c, \nu_{R,3}^c, \nu_{R,4}^c, S_1^c, S_2^c)$ is the obvious extension of the MES structure from Eq. (1) for the $3+2$ scheme, while $\mathcal{M}_S$ denotes the mass matrix of S_{3-6}, i.e.,

$$\mathcal{M}_S = \begin{pmatrix} 0 & y_1\langle\phi\rangle & 0 & 0 \\ y_1\langle\phi\rangle & 0 & 0 & 0 \\ 0 & 0 & 0 & y_2\langle\phi\rangle \\ 0 & 0 & y_2\langle\phi\rangle & 0 \end{pmatrix} , \tag{23}$$

resulting in two Dirac fermions, decoupled from the neutrino sector.

Compared to the $3+1$ scheme discussed so far, the scalar sector is identical, whereas the dark matter sector is slightly modified because we have only two stable Dirac fermions (protected by the remaining discrete gauge group $\mathbb{Z}_5$) instead of three, but two light sterile neutrinos instead of one. This does not influence the qualitative behavior significantly.

The expressions from Sec. III A for the neutrino masses go through in the same manner, i.e., we still have

$$\mathcal{M}_\nu^{3\times 3} \simeq -m_D M_R^{-1} m_D^T + m_D M_R^{-1} m_S \,(m_S^T M_R^{-1} m_S)^{-1} m_S^T M_R^{-1} m_D^T \,, \quad \mathcal{M}_{\nu_s}^{2\times 2} \simeq -m_S^T M_R^{-1} m_S \,, \tag{24}$$

for the masses, where we assumed $m_D \ll m_S \ll M_R$, and m_D, m_S and M_R are $3\times n$, $n\times 2$, and $n\times n$ matrices, respectively. The number n of gauge group singlets ν_R should be at least 4 to make enough light neutrinos massive, $n \geq 5$ makes all five light neutrinos massive. The active–sterile mixing is again $\mathcal{O}(m_D/m_S)$, so the required values $\mathcal{O}(0.1)$ put the $U(1)'$ breaking scale naturally in the TeV range.

Let us briefly comment on the thermal evolution of the universe in this model. Seeing as the number of degrees of freedom is smaller (larger) at $t_{\rm sep}$ (t_f) compared to the $3+1$ scheme of Sec. III C 2, the sterile neutrino bath is colder than the SM bath (prior to neutrino decoupling) by a factor of $\simeq 0.75$. Without active–sterile neutrino oscillations, this would mean that the two sterile neutrinos effectively only contribute $\Delta N_{\rm eff} \simeq 0.6$ additional neutrino species to the energy density, alleviating cosmological bounds. It is of course to be expected that active–sterile oscillations before neutrino decoupling generate thermal equilibrium among the neutrinos, giving rise to the usual constraints.

For completeness, we give an assignment for the $3+3$ case, which has been fitted to the neutrino anomalies in Ref. [31]. To make at least five light neutrinos massive, we need five ν_R. A possible charge assignment for nine S_j is then $(7,7,7,2,-9,-1,-6,-4,-3)$, with one scalar $\phi\sim 7$. This leads to three light sterile neutrinos and three stable Dirac DM particles (protected by the remaining $\mathbb{Z}_7$).

B. Majorana Dark Matter

Having focused on Dirac DM in the main text for no particular reason, we will now give an example with Majorana DM. For the $3+1$ MES scheme, we take the exotic charges $(6,-3,-3,2,-8,-1,7)$ for the S_i and one scalar with charge $Y'(\phi) = 6$. The VEV of ϕ breaks $U(1)' \to \mathbb{Z}_6$, S_1 will again become the sterile neutrino, while S_2 and S_3 share the most general Majorana mass matrix—which we can take to be diagonal without loss of generality—resulting in two Majorana fermions $\Psi_{1,2}$. (S_4, S_5, S_6, S_7) share the mass matrix

$$\mathcal{M}_S = \begin{pmatrix} 0 & y_1\langle\phi\rangle & 0 & 0 \\ y_1\langle\phi\rangle & 0 & 0 & 0 \\ 0 & 0 & 0 & y_2\langle\phi\rangle \\ 0 & 0 & y_2\langle\phi\rangle & 0 \end{pmatrix} , \tag{25}$$

resulting in two Dirac fermions $\Psi_{3,4}$; all Ψ_j are decoupled from the neutrino sector. These particles form representations $\Psi_{1,2}\sim 3\sim(1,0)$, $\Psi_3\sim 2\sim(0,2)$, $\Psi_4\sim 1\sim(1,1)$ under $\mathbb{Z}_6\cong\mathbb{Z}_2\times\mathbb{Z}_3$, so depending on the mass spectrum, we can obtain a stable Majorana fermion.

C. Unstable Dark Matter

The charges for the S_i and ϕ discussed so far have been chosen in such a way that the spontaneous breaking of $U(1)'$ leaves a nontrivial $\mathbb{Z}_N$ that stabilizes the DM candidates. This is of course not a generic feature of exotic charges, but just a convenient choice to obtain exactly stable particles. Let us briefly comment on unstable DM candidates: Taking $(1,-10,9,-7,6,-11,12)$ for the $3+1$ scheme with a scalar $\phi\sim 1$ gives three Dirac DM candidates $\Psi_1 = S_2 + S_3^c$, $\Psi_2 = S_4 + S_5^c$, $\Psi_3 = S_6 + S_7^c$, which are independently stable due to an accidental global $U(1)^3$ symmetry. However, with this charge assignment, there is no leftover $\mathbb{Z}_N$ symmetry protecting this stability. Similar to the discussion in Sec. III C 1, we can discuss higher-dimensional operators. For the charge assignment here, there are already dimension-five operators

$$\phi^2 \overline{S}_3 S_4^c/\Lambda\,, \qquad \phi^2 \overline{S}_3^c S_6/\Lambda\,, \qquad \phi^2 \overline{S}_2 S_7^c/\Lambda\,, \tag{26}$$

which break the global $U(1)^3$ to a $U(1)$ symmetry, so only one stable Dirac fermion survives. Even this stability is not exact, as there are operators like $\phi^6 \overline{S}_5^c \nu_R/\Lambda^5$ which break the global $U(1)$ and lead to DM decay. In this particular example—and for $\Lambda \sim M_{\rm Pl}$—the decay would be suppressed enough to still allow valid DM, but in principle there are charge assignments with decaying DM, or even no DM candidate at all.

V. CONCLUSION

Generating small sterile neutrino masses via the same seesaw mechanism that suppresses active neutrino masses requires a specific structure in the neutral fermion mass matrix. We showed how this so-called MES structure can be obtained from a new spontaneously broken $U(1)'$ symmetry, under which the "sterile" neutrino is charged. Heavily mixed eV-scale steriles hint at a $U(1)'$ breaking scale around TeV. Additional anomaly-cancelling fermions need to carry exotic $U(1)'$ charges in order to not spoil the MES structure, which coincidentally stabilizes one or more of them (without the need for additional discrete symmetries). The main connection between this multicomponent dark matter sector and the Standard Model is the active–sterile mixing (neutrino portal). We discussed how the dark matter annihilation almost exclusively into sterile neutrinos can be used to obtain the measured relic density, and also the interplay with the other two portals (Higgs- and kinetic-mixing portals).

We focussed on a few specific examples, but the presented framework of exotic charges obviously provides a rich playground for model building, depending on the used charges and number of new particles. Worthwhile extensions with $U(1)'$ charged SM fermions, e.g. $B-L$ type symmetries, can be obtained with slightly more complicated scalar sectors and will be discussed elsewhere.[††]

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