

A model of radiative neutrino mass: with or without dark matter

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ABSTRACT: We present a three-loop model of neutrino mass whose most-general Lagrangian possesses a softly-broken accidental Z_2 symmetry. In the limit that a single parameter vanishes, $\lambda \rightarrow 0$, the Z_2 symmetry becomes exact and the model contains a stable dark-matter candidate. However, even for finite $\lambda \ll 1$, long-lived dark matter is possible, giving a unified solution to the neutrino mass and dark matter problems that does not invoke a new symmetry. Taken purely as a neutrino mass model, the new physics can be at the TeV scale. When dark matter is incorporated, however, only a singlet scalar can remain this light, though the dark matter can be tested in direct-detection experiments.

KEYWORDS: Beyond Standard Model, Cosmology of Theories beyond the SM, Neutrino Physics, Phenomenological Models

ARXIV EPRINT: [1404.5917](https://arxiv.org/abs/1404.5917)

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1 Introduction

The observation of neutrino oscillations in solar, atmospheric and reactor experiments confirms that neutrinos are massive and that the Standard Model (SM) is incomplete. Another strong motivation for beyond-SM physics comes from astrophysical observations, which motivate a new gravitating particle species referred to as dark matter (DM). It is sensible to ask if these two problems could have a common solution. Models with radiative neutrino mass [1–3] offer a promising direction for a unified solution to these problems (for a discussion of radiative models see e.g. [4]). If the coupling to DM is related to the source of lepton number symmetry breaking, DM can propagate inside the loop diagram that generates neutrino mass, killing the proverbial two birds with one stone.

An early proposal along these lines was put forward by Krauss, Nasri and Trodden (KNT) [6] (for analysis see refs. [7–10]). In this paper we investigate a three-loop model of neutrino mass that is related to the KNT model. The model differs in its field content but employs a three-loop diagram with the same topology. The use of distinct beyond-SM multiplets produces some key differences. Recall that the KNT model utilizes a discrete (Z_2) symmetry, which serves two purposes: it precludes tree-level neutrino mass, which

would otherwise dominate the loop mass, and it gives a stable particle that is taken as the DM. Consequently, the KNT model gives a unified solution to the neutrino mass and DM problems.

Different from the KNT model, the present model does not require a new symmetry to preclude tree-level neutrino mass, despite sharing the same loop-topology. It is therefore a viable model of radiative neutrino mass independent of any DM considerations. Interestingly, the most-general Lagrangian for the model possesses a softly-broken accidental Z_2 symmetry. In the limit where a single parameter vanishes, $\lambda \rightarrow 0$, this symmetry becomes exact and the model contains a stable DM candidate. As a result, the DM width goes like $\Gamma_{DM} \propto \lambda^2$ for nonzero λ , and one can always make this sufficiently small to obtain long-lived DM, or simply take $\lambda \rightarrow 0$ for absolutely stable DM. Thus, DM is possible with or without the Z_2 symmetry. This gives a unified solution to the DM and neutrino mass problems that does not require a new symmetry. Importantly, the limit $\lambda \rightarrow 0$ does not affect the predictions for neutrino mass. The Z_2 symmetry is essentially the same one found in the KNT model (and the related triplet model [11]), though in those cases the most-general Lagrangian contains multiple symmetry breaking terms, including ones that give tree-level neutrino mass.

We shall see that the phenomenology of the model depends on the region of parameter space considered. Taken purely as a model of neutrino mass, the new physics can be at the TeV scale and may be probed in collider experiments. When DM is incorporated one requires $M_{DM} \sim 10$ TeV, putting some of the new multiplets beyond the reach of colliders. None the less, a singly-charged scalar that appears in the model can remain at the TeV scale, with or without the inclusion of DM. Even when DM is included, prospects for testing the model in direct-detection experiments are good.

We note that one-loop models of neutrino mass that admit DM candidates but do not require a symmetry to exclude tree-level masses exist [12], with one model further studied in ref. [13]. Other works studying connections between neutrino mass and DM include refs. [14–39]. For a review see [40].

The layout of this paper is as follows. In section 2 we describe the basic details of the model. Neutrino masses are calculated in section 3 and important flavor-changing constraints are discussed in section 4. We consider DM in section 5, discussing the issue of longevity and the relic abundance. Our main numerical results and discussion are given in section 6 and we comment on collider phenomenology in section 7. We briefly describe interesting generalizations of our model in section 8, and conclude in section 9.

2 The model

We extend the SM to include a charged scalar singlet, $S^+ \sim (1, 1, 2)$, a complex scalar quintuplet, $\phi \sim (1, 5, 2)$, and a real fermion quintuplet, $\mathcal{F} \sim (1, 5, 0)$. We write the exotics in symmetric-matrix form as ϕ_{abcd} and \mathcal{F}_{abcd} , where

$$\phi_{1111} = \phi^{+++}, \quad \phi_{1112} = \frac{\phi^{++}}{\sqrt{4}}, \quad \phi_{1122} = \frac{\phi^+}{\sqrt{6}}, \quad \phi_{1222} = \frac{\phi^0}{\sqrt{4}}, \quad \phi_{2222} = \phi^-, \quad (2.1)$$

$$\mathcal{F}_{1111} = \mathcal{F}_L^{++}, \quad \mathcal{F}_{1112} = \frac{\mathcal{F}_L^+}{\sqrt{4}}, \quad \mathcal{F}_{1122} = \frac{\mathcal{F}_L^0}{\sqrt{6}}, \quad \mathcal{F}_{1222} = \frac{(\mathcal{F}_R^+)^c}{\sqrt{4}}, \quad \mathcal{F}_{2222} = (\mathcal{F}_R^{++})^c.$$

Note that ϕ^+ and ϕ^- are distinct fields and, in particular, $\phi^- \neq (\phi^+)^*$. The Lagrangian for the model contains the terms

$$\mathcal{L} \supset \mathcal{L}_{\text{SM}} + \{f_{\alpha\beta} \overline{L}_\alpha^c L_\beta S^+ + g_{i\alpha} \overline{\mathcal{F}}_i \phi e_{\alpha R} + \text{H.c.}\} - \frac{1}{2} \overline{\mathcal{F}}_i^c \mathcal{M}_{ij} \mathcal{F}_j - V(H, S, \phi). \quad (2.2)$$

We label lepton flavors by lower-case Greek letters, $\alpha, \beta \in \{e, \mu, \tau\}$, while exotic fermion generations are labeled by i . The superscript ‘‘c’’ denotes charge conjugation. The Lagrangian shows that the multiplets ϕ and \mathcal{F} are sequestered from the SM neutrinos. None the less, they play a key role in enabling neutrino mass, as we shall shortly see.

The explicit expansion of the fermion mass term gives

$$\begin{aligned} & -\frac{1}{2} (\overline{\mathcal{F}}_i^c)_{abcd} \mathcal{M}_{ij} (\mathcal{F}_j)_{efgh} \epsilon^{ae} \epsilon^{bf} \epsilon^{cg} \epsilon^{dh} + \text{H.c.} \\ & = -\overline{\mathcal{F}}_{iR}^{++} \mathcal{M}_{ij} \mathcal{F}_{jL}^{++} + \overline{\mathcal{F}}_{iR}^+ \mathcal{M}_{ij} \mathcal{F}_{jL}^+ - \frac{1}{2} (\overline{\mathcal{F}}_{iL}^0)^c \mathcal{M}_{ij} \mathcal{F}_{jL}^0 + \text{H.c.} \\ & = -\overline{\mathcal{F}}_i^{++} \mathcal{M}_{ij} \mathcal{F}_j^{++} - \overline{\mathcal{F}}_i^+ \mathcal{M}_{ij} \mathcal{F}_j^+ - \frac{1}{2} \overline{\mathcal{F}}_i^0 \mathcal{M}_{ij} \mathcal{F}_j^0, \end{aligned} \quad (2.3)$$

where, in the last line, we define:

$$\mathcal{F}^{++} = \mathcal{F}_L^{++} + \mathcal{F}_R^{++}, \quad \mathcal{F}^+ = \mathcal{F}_L^+ - \mathcal{F}_R^+, \quad \mathcal{F}^0 = \mathcal{F}_L^0 + (\mathcal{F}_L^0)^c. \quad (2.4)$$

Here \mathcal{F}^0 is clearly a Majorana fermion, while the other four components of \mathcal{F} partner-up to give two charged (Dirac) fermions. Without loss of generality we work in a basis with $\mathcal{M}_{ij} = \text{diag}(M_1, M_2, M_3)$, where the masses are ordered as $M_1 < M_2 < M_3$. In what follows, we will use $M_F \equiv M_1$ for the DM mass.

In terms of these fields the Yukawa couplings involving the new fermions are written as

$$\begin{aligned} g_{i\alpha} (\overline{\mathcal{F}}_i)^{abcd} \phi_{abcd} e_{\alpha R} &= g_{i\alpha} \left\{ \phi^{+++} \overline{\mathcal{F}}_i^{+++} P_R e_\alpha + \phi^{++} \overline{\mathcal{F}}_i^{++} P_R e_\alpha \right. \\ & \left. + \phi^+ \overline{\mathcal{F}}_i^0 P_R e_\alpha - \phi^0 (\overline{e_\alpha})^c P_R \mathcal{F}_i^+ + \phi^- (\overline{e_\alpha})^c P_R \mathcal{F}_i^{++} \right\}, \end{aligned} \quad (2.5)$$

where P_R is a standard projection operator. The extra minus sign is due to the definition of \mathcal{F}^+ .

We consider the parameter space where the SM Higgs breaks the electroweak symmetry via the nonzero vacuum value $\langle H \rangle \neq 0$, while $\langle \phi \rangle = 0$, so the SM tree-level value of the ρ -parameter is not modified. Before turning to neutrino mass we would like to discuss a few features of the model. To this end, let us briefly consider the theory in the absence of the singlet S . In this case the scalar potential is

$$V(H, \phi) = V(H) + V(\phi) + V_m(H, \phi), \quad (2.6)$$

where the mixing potential is

$$V_m(H, \phi) = \lambda_{H\phi 1} (\phi^*)^{abcd} \phi_{abcd} (H^*)^e H_e + \lambda_{H\phi 2} (\phi^*)^{abcd} \phi_{ebcd} (H^*)^e H_a. \quad (2.7)$$

This potential $V(H, \phi)$ possesses an accidental $U(1)$ symmetry, $\phi \rightarrow e^{i\theta}\phi$. However, the coupling to \mathcal{F} breaks this symmetry to a discrete subgroup, due to the Majorana mass. Therefore in the absence of S , the theory has an accidental Z_2 symmetry:

$$\{\phi, \mathcal{F}\} \rightarrow \{-\phi, -\mathcal{F}\}. \tag{2.8}$$

Adding S to the theory, the full potential can be written as

$$V(H, S, \phi) = V(H, \phi) + V(S) + V_m(S, \phi) + V_m(H, S) + V_m(H, S, \phi), \tag{2.9}$$

where the first four terms in this potential all preserve the discrete symmetry, and the last mixing-potential is given by

$$V_m(H, S, \phi) = \frac{\lambda_S}{4}(S^-)^2 \phi_{abcd}\phi_{efgh}\epsilon^{ae}\epsilon^{bf}\epsilon^{cg}\epsilon^{dh} + \lambda S^-(\phi^*)^{abcd}\phi_{abef}\phi_{cdjl}\epsilon^{ej}\epsilon^{fl} + \text{H.c.} \tag{2.10}$$

The first term in this potential also preserves the Z_2 symmetry, leaving the second term as the sole source of Z_2 symmetry-breaking in the full theory. Thus, in the limit $\lambda \rightarrow 0$ the theory possesses the Z_2 symmetry $\{\phi, \mathcal{F}\} \rightarrow \{-\phi, -\mathcal{F}\}$, making $\lambda \ll 1$ technically natural. This symmetry is analogous to that invoked in both the KNT model [6] and the three-loop model with triplets [11]. In the limit that $\lambda \rightarrow 0$ a stable particle emerges, which we return to in section 5.

If the Z_2 symmetry were exact, it would prevent mixing between \mathcal{F} and the SM leptons. Consequently any such mixing must be generated radiatively and must involve the coupling λ . This mixing is of a sufficiently high order as to be negligible, though to be certain one can always choose λ sufficiently small to make the mixing negligible. We can therefore ignore any mixing between \mathcal{F} and the SM.

At tree-level the components of \mathcal{F} are mass-degenerate, while the components of ϕ experience a mild splitting due to the $\lambda_{H\phi 2}$ -term in $V_m(H, \phi)$. For $M_\phi \gtrsim \mathcal{O}(\text{TeV})$ this mass-splitting is not significant and is negligible for $\lambda_{H\phi 2} \lesssim 0.1$. Thus, to good approximation the components of \mathcal{F} and ϕ are degenerate at tree-level, with masses $M_{\mathcal{F}}$ and M_ϕ , respectively. Radiative corrections lift these mass degeneracies. For example, loops involving SM gauge bosons induce splittings of $M_{\mathcal{F}^{++}} - M_{\mathcal{F}^+} \simeq 490 \text{ MeV}$, and $M_{\mathcal{F}^+} - M_{\mathcal{F}^0} \simeq 163 \text{ MeV}$, among the components of \mathcal{F} , leaving \mathcal{F}^0 as the lightest state once loop-corrections are incorporated [41–43]. Similar splittings are induced for the components of ϕ [41]. For most purposes these small splittings can be neglected.

We note that the fermions $\mathcal{F} \sim (1, 5, 0)$ employed in this model were studied in a number of other contexts. They allow a generalization of the Type-III seesaw mechanism [44] that achieves neutrino mass via a low-energy effective operator of mass-dimension $d = 9$ [42, 45, 46]. Similarly they permit a generalized inverse seesaw mechanism [48]. The neutral component of the fermion is also the favored ‘‘Minimal DM’’ candidate [41]. For related phenomenological studies see refs. [49, 50].

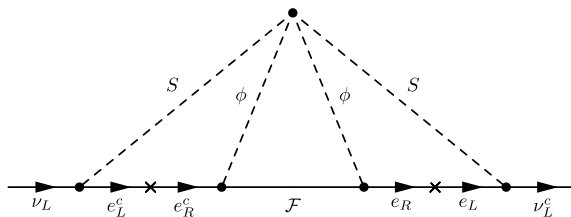


Figure 1. Three-loop diagram for radiative neutrino mass, where S and ϕ are new scalars and \mathcal{F} is an exotic fermion.

3 Three-loop radiative neutrino masses

The Yukawa Lagrangian is not sufficient to break lepton number symmetry. However, as just mentioned, the scalar potential contains the terms

$$\begin{aligned} V(H, S, \phi) &\supset \frac{\lambda_S}{4} (S^-)^2 \phi_{abcd} \phi_{efgh} \epsilon^{ae} \epsilon^{bf} \epsilon^{cg} \epsilon^{dh} + \frac{\lambda_S^*}{4} (S^+)^2 (\phi^*)^{abcd} (\phi^*)^{efgh} \epsilon_{ae} \epsilon_{bf} \epsilon_{cg} \epsilon_{dh} \\ &= \frac{\lambda_S}{2} (S^-)^2 \left\{ \phi^{+++} \phi^- - \phi^{++} \phi^0 + \frac{1}{2} \phi^+ \phi^+ \right\} + \text{H.c.} \end{aligned} \quad (3.1)$$

When combined with the Yukawa couplings, these ensure that lepton number symmetry is explicitly broken in the model. Consequently, Majorana neutrino masses are generated radiatively, appearing at the three-loop level as shown in figure 1. There are actually five distinct diagrams, corresponding to the sets $\{\phi^+, \mathcal{F}^0, (\phi^+)^*\}$, $\{\phi^{++}, (\mathcal{F}^+)^c, \phi^0\}$, $\{\phi^0, \mathcal{F}^+, \phi^{--}\}$, $\{\phi^{+++}, (\mathcal{F}^+)^c, (\phi^-)^*\}$ and $\{\phi^-, \mathcal{F}^{++}, \phi^{---}\}$ propagating in the inner loop in figure 1.

In the limit where the mass-splitting among components of ϕ and \mathcal{F} are neglected, the loop-diagrams gives

$$(\mathcal{M}_\nu)_{\alpha\beta} = \frac{5\lambda_S}{(4\pi^2)^3} \frac{m_\gamma m_\delta}{M_\phi} f_{\alpha\gamma} f_{\beta\delta} g_{\gamma i}^* g_{\delta i}^* \times F \left(\frac{M_i^2}{M_\phi^2}, \frac{M_S^2}{M_\phi^2} \right). \quad (3.2)$$

Here the function F encodes the loop integrals (see appendix A) and has the same form as given in ref. [9]. M_S is the charged-singlet mass and M_ϕ is the mass of the degenerate members of ϕ .

The entries in the neutrino mass matrix $(\mathcal{M}_\nu)_{\alpha\beta}$ may be related to the mass eigenvalues and the elements of the Pontecorvo-Maki-Nakawaga-Sakata (PMNS) mixing matrix [51, 52]:

$$(\mathcal{M}_\nu)_{\alpha\beta} = [U_\nu \cdot \text{diag}(m_1, m_2, m_3) \cdot U_\nu^\dagger]_{\alpha\beta}. \quad (3.3)$$

We parameterize the PMNS matrix as

$$U_\nu = \begin{pmatrix} c_{12}c_{13} & c_{13}s_{12} & s_{13}e^{-i\delta_D} \\ -c_{23}s_{12} - c_{12}s_{13}s_{23}e^{i\delta_D} & c_{12}c_{23} - s_{12}s_{13}s_{23}e^{i\delta_D} & c_{13}s_{23} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta_D} & -c_{12}s_{23} - c_{23}s_{12}s_{13}e^{i\delta_D} & c_{13}c_{23} \end{pmatrix} \times U_p, \quad (3.4)$$

where the Majorana phases $\theta_{\alpha,\beta}$ appear in the matrix $U_p = \text{diag}(1, e^{i\theta_{\alpha/2}}, e^{i\theta_{\beta/2}})$, and δ_D is the Dirac phase. The dependence on the mixing angles is denoted by $s_{ij} \equiv \sin \theta_{ij}$, $c_{ij} \equiv$

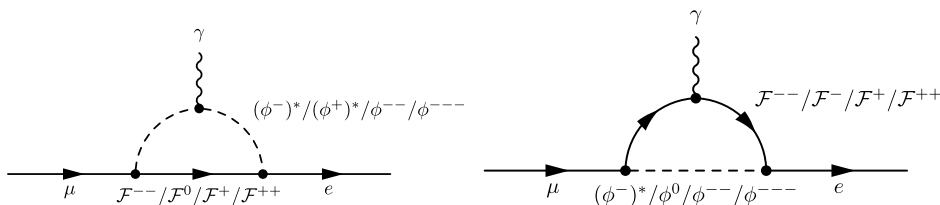


Figure 2. Diagrams for $\mu \rightarrow e + \gamma$ due to the Z_2 -odd fields $\mathcal{F} \sim (1, 5, 0)$ and $\phi \sim (1, 5, 2)$. There is also a similar diagram involving the scalar $S \sim (1, 1, 2)$.

$\cos \theta_{ij}$. Analysis of neutrino experimental data gives the best-fit values for the mass-squared differences and mixing angles are $s_{13}^2 = 0.025_{-0.003}^{+0.003}$, $s_{23}^2 = 0.43_{-0.03}^{+0.03}$, $s_{12}^2 = 0.320_{-0.017}^{+0.016}$, $|\Delta m_{13}^2| = 2.55_{-0.09}^{+0.06} \times 10^{-3} \text{eV}^2$, and $\Delta m_{21}^2 = 7.62_{-0.19}^{+0.19} \times 10^{-5} \text{eV}^2$ [53, 54]. Matching to these values determines the regions of parameter space with viable neutrino masses.

4 Experimental constraints

The Yukawa couplings $g_{i\alpha}$ induce flavor changing processes like $\mu \rightarrow e + \gamma$. At the one-loop level, there are two classes of diagrams containing \mathcal{F} and ϕ that one should consider, as shown in figure 2. Note, however, that diagrams with the photon attached to the internal fermion come in pairs which differ by an overall sign. The coherent sum of the corresponding amplitudes vanishes in the limit that the small mass-splitting are neglected.¹ A similar cancelation occurs between the diagrams containing singly-charged scalars in figure 2a. Calculating the diagrams in figure 2, and adding the diagram involving the singlet S , one finds that the branching fraction for $\mu \rightarrow e + \gamma$ is given by

$$\begin{aligned}
 B(\mu \rightarrow e\gamma) &= \frac{\Gamma(\mu \rightarrow e + \gamma)}{\Gamma(\mu \rightarrow e + \nu + \bar{\nu})} \\
 &\simeq \frac{\alpha v^4}{384\pi} \times \left\{ \frac{|f_{\mu\tau} f_{\tau e}^*|^2}{M_S^4} + \frac{900}{M_\phi^4} \left| \sum_i g_{ie}^* g_{i\mu} F_2(M_i^2/M_\phi^2) \right|^2 \right\}. \quad (4.1)
 \end{aligned}$$

Here the function $F_2(R) = (R - 1)^{-4} [1 - 6R + 3R^2 + 2R^3 - 6R^2 \log R]/6$ is a standard one-loop function. A simple change of the flavor labels in eq. (4.1) allows one to obtain the related expression for $B(\tau \rightarrow \mu + \gamma)$.

When the final-state electrons in figure 2 are replaced with muons, the diagrams contribute to the magnetic moment of the muon. Similar arguments to those just used also apply for the calculation of the magnetic moment; for example, the diagrams with the photon attached to the internal-fermion line cancel in the limit the mass-splitting is neglected. The remaining diagrams give

$$|\delta a_\mu| = \frac{m_\mu^2}{16\pi^2} \left\{ \sum_{\alpha \neq \mu} \frac{|f_{\mu\alpha}|^2}{6M_S^2} + \sum_i \frac{5|g_{i\mu}|^2}{M_\phi^2} F_2(M_i^2/M_\phi^2) \right\}, \quad (4.2)$$

where once again a diagram involving the scalar S must be included.

¹Said differently, the cancelation occurs because $\sum_{\mathcal{F}} Q_{\mathcal{F}} = 0$ for all non-trivial $SU(2)$ multiplets with vanishing hypercharge ($Q_{\mathcal{F}}$ are the charges of the components of \mathcal{F}).

Null-results from searches for neutrino-less double-beta decay provide an additional constraint of $(\mathcal{M}_\nu)_{ee} \lesssim 0.35$ eV [55] that must also be considered. We find that this constraint is easily satisfied in the model. Next generation experiments will improve this bound to the level of $(\mathcal{M}_\nu)_{ee} \lesssim 0.01$ eV [56, 57].

5 Dark matter

5.1 Dark matter longevity

In the preceding we considered the most-general version of the model, where all parameters allowed by the gauge symmetries are included. We showed that the model can generate viable neutrino masses for a wide range of exotic mass scales. In this section we turn our attention to DM, to determine whether the model can provide a unified solution to the DM and neutrino mass problems. The first issue to discuss is the matter of DM longevity. As noted earlier, the model possesses a softly broken Z_2 symmetry, $\{\phi, \mathcal{F}\} \rightarrow \{-\phi, -\mathcal{F}\}$, which becomes exact in the limit $\lambda \rightarrow 0$. This suggests that the model could also give a dark matter candidate.

There are two candidates for the DM in this model, either the scalar ϕ^0 or the lightest neutral fermion \mathcal{F}_1^0 . However, ϕ^0 couples to the Z boson and can be excluded by direct-detection experiments, due to tree-level interactions with SM matter and an absence of any splitting between the real and imaginary components of ϕ^0 . This leaves \mathcal{F}_1^0 as the sole candidate, suggesting Majorana DM and requiring $M_{DM} = M_F < M_\phi$. Consider the case with $\lambda \neq 0$. Then, there are two types of one-loop \mathcal{F}_1^0 decays that can dominate, depending on the ordering of M_S and M_{DM} , namely

$$\begin{aligned} \mathcal{F}_1^0 &\longrightarrow S + 3e && \text{for} && M_S < M_F, \\ \mathcal{F}_1^0 &\longrightarrow 4e + \nu && \text{for} && M_F < M_S. \end{aligned} \tag{5.1}$$

The corresponding widths are approximately

$$\begin{aligned} \Gamma(\mathcal{F}_1^0 \rightarrow S + 3e) &\sim |\lambda|^2 M_F \frac{|g_{1\alpha} g_j \beta g_{j\gamma}|^2}{(16\pi^2)^2} \Phi_{4\text{-body}} && \text{for} && M_S \ll M_F, \\ \Gamma(\mathcal{F}_1^0 \rightarrow 4e + \nu) &\sim |\lambda|^2 M_F \frac{|g_{1\alpha} g_j \beta g_{j\gamma} f_{\delta\epsilon}|^2}{(16\pi^2)^2} \Phi_{5\text{-body}} && \text{for} && M_F < M_S, \end{aligned} \tag{5.2}$$

where $\Phi_{n\text{-body}}$ denotes the n -body phase space factor. Due to the presence of the softly-broken accidental Z_2 symmetry, one can always choose nonzero $\lambda \ll 1$ sufficiently small to ensure adequate dark-matter longevity. This provides a simple way to include a DM candidate without recourse to an additional symmetry. The limit $\lambda \rightarrow 0$ then smoothly interpolates to the Z_2 -symmetric case, making \mathcal{F}_1^0 absolutely stable. Importantly, neutrino masses are not sensitive to this limit, and viable masses can be obtained irrespective of DM considerations.

5.2 Relic density

Taking the neutral fermion \mathcal{F}_1^0 as the DM candidate, there are two classes of interactions that maintain thermal contact with the SM in the early universe. This includes processes

mediated by $SU(2)_L$ gauge bosons, which can be calculated in the $SU(2)$ -symmetric limit, and others mediated by the scalar ϕ . One must also include coannihilation processes in the calculation, due to the small mass-splitting between the charged and neutral fermions.

The annihilation of DM due to ϕ -exchange give

$$\sigma(2\mathcal{F}^0 \rightarrow \ell_\beta^+ \ell_\alpha^-) \times v_r = \frac{|g_{1\alpha}^* g_{1\beta}|^2 M_F^2 (M_F^4 + M_\phi^4)}{48\pi (M_F^2 + M_\phi^2)^4} \times v_r^2 \equiv \sigma_{0,0}^{\alpha\beta} \times v_r, \quad (5.3)$$

where $v_r = 2v$ is the relative velocity of the dark matter in the centre-of-mass frame. There are no s -wave annihilations in this expression as the DM is a Majorana fermion and we neglected final-state lepton masses. There are no coannihilations mediated by ϕ , though one must include the annihilations for singly charged fermions:

$$\sigma(\mathcal{F}^- \mathcal{F}^+ \rightarrow \ell_\beta^+ \ell_\alpha^-) \times v_r = \frac{|g_{1\alpha}^* g_{1\beta}|^2 M_F^2 (M_F^4 + M_\phi^4)}{48\pi (M_F^2 + M_\phi^2)^4} \times v_r^2 \equiv \sigma_{-,+}^{\alpha\beta} \times v_r, \quad (5.4)$$

and doubly-charged fermions

$$\sigma(\mathcal{F}^{--} \mathcal{F}^{++} \rightarrow \ell_\beta^+ \ell_\alpha^-) \times v_r = \frac{|g_{1\alpha}^* g_{1\beta}|^2 M_F^2 (M_F^4 + M_\phi^4)}{48\pi (M_F^2 + M_\phi^2)^4} \times v_r^2 \equiv \sigma_{--,++}^{\alpha\beta} \times v_r. \quad (5.5)$$

For annihilations and coannihilations involving $SU(2)_L$ gauge bosons we can make use of known results in the literature [58, 59].

In the limit where the mass-splitting between fermion components vanishes, $\Delta M_{\mathcal{F}} \rightarrow 0$, we add annihilation and coannihilation channels together with the standard method [60] to obtain

$$\sigma_{\text{eff}}(2\mathcal{F} \rightarrow SM) \times v_r = \frac{1}{g_{\text{eff}}^2} \left[\sigma_W \times v_r + \sum_{\alpha,\beta} \left\{ g_0^2 \sigma_{0,0}^{\alpha\beta} + 2g_\pm \sigma_{-,+}^{\alpha\beta} + 2g_{\pm\pm} \sigma_{--,++}^{\alpha\beta} \right\} \times v_r \right], \quad (5.6)$$

where the $SU(2)_L$ channels are denoted by

$$\sigma_W \equiv \frac{\pi \alpha_2^2}{2M_F^2 v_r} \left\{ 2070 + \frac{1215}{2} v_r^2 \right\}, \quad (5.7)$$

and $g_{\text{eff}} = g_0 + 2g_\pm + g_{\pm\pm}$, with $g_0 = g_\pm = g_{\pm\pm} = 2$. The ϕ -exchange cross sections are defined above.

5.3 Direct detection

The DM candidate does not couple to quarks at tree-level due to its vanishing hypercharge and, being a Majorana fermion, there are no radiative magnetic-dipole interactions with SM gauge bosons. Exchange of W bosons generates the three one-loop diagrams in figure 3, which are relevant for direct-detection experiments. The scattering contains both spin-dependent and spin-independent contributions. However, the former are suppressed by the DM mass, expected to be $M_F \sim 10$ TeV in our case, giving highly-suppressed spin-dependent scattering cross sections.

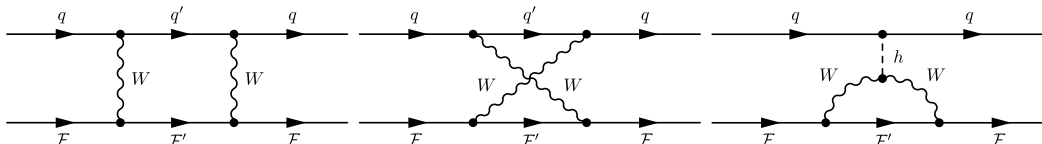


Figure 3. Feynman diagrams for direct-detection experiments.

The dominant interaction is therefore spin-independent scattering, with a cross section determined by SM interactions:

$$\sigma_{\text{SI}}(\mathcal{F}^0 N \rightarrow \mathcal{F}^0 N) \simeq \frac{9\pi\alpha_2^4 M_A^4 f^2}{M_W^2} \left[\frac{1}{M_W^2} + \frac{1}{M_h^2} \right]^2. \quad (5.8)$$

Here the DM scatters off a target nucleus A with mass M_A and we use a standard parameterizations for the nucleon,

$$\langle N | \sum_q m_q \bar{q}q | N \rangle = f m_N, \quad (5.9)$$

with m_N being the nucleon mass. We use $f \approx 1/3$, though this is subject to the standard QCD uncertainties. The resulting cross section per nucleon is of order $\sigma_{\text{SI}} \simeq 10^{-46} \text{cm}^2$, which is beyond the current sensitivity of LUX [61], but within reach of forthcoming experiments like SuperCDMS [62]. Discovery prospects are therefore promising.

6 Results and discussion

Using the results from the preceding sections, we can determine the parameter space where viable neutrino masses are obtained and the correct DM relic-density is realized. Here, we present the results from our numerical scans of the parameter space. We find that neutrino masses can be obtained for a range of parameter space, including the fermion and scalar masses. It also appears that the observed DM relic abundance can be generated. Whenever we consider \mathcal{F}_1 as DM, we assume λ is sufficiently small to ensure DM longevity. In our numerical scan, we consider the following range for the model parameters

$$\begin{aligned} |f_{\alpha\beta}|^2, |g_{i\alpha}|^2 &\lesssim 9, & 500 \text{ GeV} &\leq M_F \leq 10 \text{ TeV}, \\ 300 \text{ GeV} &\leq M_S \leq 1 \text{ TeV}, & M_{2,3}, M_\phi &\gtrsim M_F, \end{aligned}$$

and we impose the constraints from neutrino mass and mixing, LFV processes and muon anomalous magnetic moment, with and without the DM relic abundance constraints. In figures 4, 5, 6 and 7, the red (blue) benchmarks represent the sets of model parameters that satisfy the constraints without (with) the DM relic density abundance.

In figure 4 we plot the relic density, $\Omega_{DM} h^2$, versus the (scaled) DM mass, $x_f = M_F/T_f$, where T_f is the freeze-out temperature. The blue points correctly reproduce the observed relic-density [63]. Viable neutrino masses are obtained for all points shown and the various constraints are satisfied; regions of parameter space that do not give the observed relic-abundance still allow a viable model of neutrino mass. For parameter space where the DM

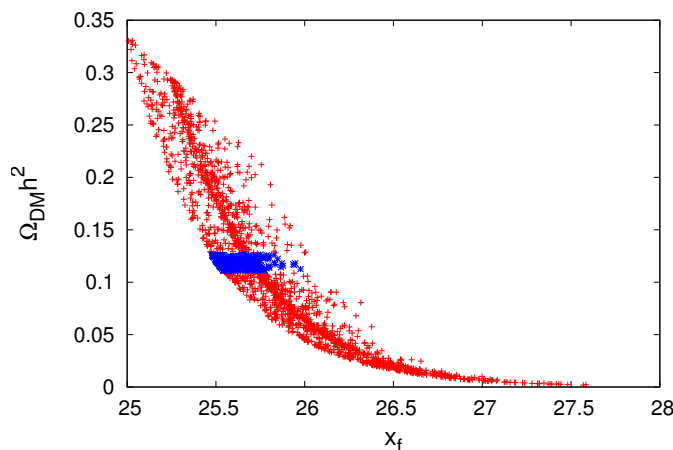


Figure 4. The DM relic abundance versus the (scaled) DM mass, $x_f = M_F/T_f$, where the blue benchmarks correspond to the physical observed value.

abundance is too large one must take λ adequately large to allow \mathcal{F}_1 to decay to the SM. In other cases one can consider small finite values of λ or simply take $\lambda \rightarrow 0$ to achieve the Z_2 -symmetric limit.

The corresponding masses for the exotic fields are shown in figure 5. In the limit where the annihilations involving ϕ are switched off, $g_{i\alpha} \rightarrow 0$, the green line in figure ($M_F = 5.844$ TeV) 5-left corresponds to the current best-fit value for the DM relic density, $\Omega_{DM} h^2 = 0.1187$. We observe that ϕ -exchange slightly modifies the value of M_F by a ratio between $[-3.3\%, 19\%]$ with the mass $M_F \sim 6$ TeV generically expected. When DM is incorporated, the fermions \mathcal{F} and the scalar ϕ are both well-beyond the reach of collider experiments. On the other hand, taken purely as a model of neutrino mass, these exotics can have $\mathcal{O}(\text{TeV})$ masses as seen in figure 5. In either case, the singlet scalar S can remain relatively light with mass $M_S = \mathcal{O}(100)$ GeV, so collider experiments should provide additional tests on the model. In our analysis we restricted our scans to parameter space with $M_F < M_\phi$, as required when \mathcal{F}_1 is the DM. When only neutrino masses are considered one could consider alternative mass orderings for the exotics, with $M_\phi < M_{\mathcal{F}}$ also possible.

The viable parameter space for the Yukawa couplings $f_{\alpha\beta}$ and $g_{i\alpha}$ is shown in figure 6. In our numerical scans we restricted these couplings to the perturbative range, $|f_{\alpha\beta}|^2, |g_{i\alpha}|^2 \lesssim 4\pi$. A reasonable spread of values are possible for $f_{\alpha\beta}$, though the scans generically require $g_{i\alpha} = \mathcal{O}(1)$. The corresponding branching fractions for the flavor-changing decays appear in figure 7. The bound on $\tau \rightarrow \mu + \gamma$ is easily satisfied, though the constraint of $B(\mu \rightarrow e + \gamma) < 5.7 \times 10^{-13}$ [64] makes the parameter space very constrained. An order of magnitude improvement in the bound on $B(\mu \rightarrow e + \gamma)$ would exclude the vast majority of the viable parameter space found in the scans. It is worth noting that with only two generations of fermions \mathcal{F}_i ($g_{3\alpha} = 0$, the bound on $B(\mu \rightarrow e + \gamma)$ is violated. Three generations of \mathcal{F}_i are therefore required to obtain agreement with constraints from lepton flavor violating processes.

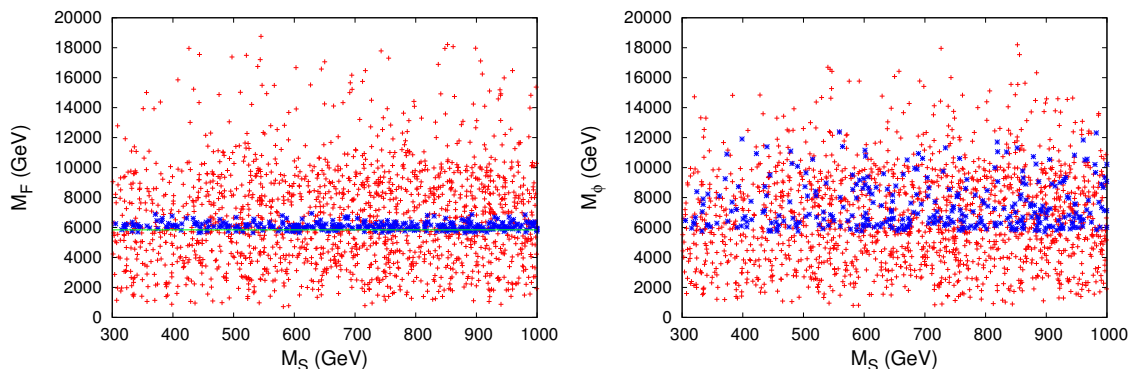


Figure 5. Allowed mass values. These achieve viable neutrino mass/mixings while satisfying the constraints. The blue points give the DM relic abundance in accordance with figure 4. Left: the lightest neutral-fermion mass versus the singlet scalar mass. When the correct relic abundance is achieved, \mathcal{F}_1 is the DM, with M_F . The green line gives the best-fit value for $\Omega_{DM}h^2$ when $g_{i\alpha} \rightarrow 0$. Right: the corresponding scalar masses, with $M_F < M_\phi$ assumed.

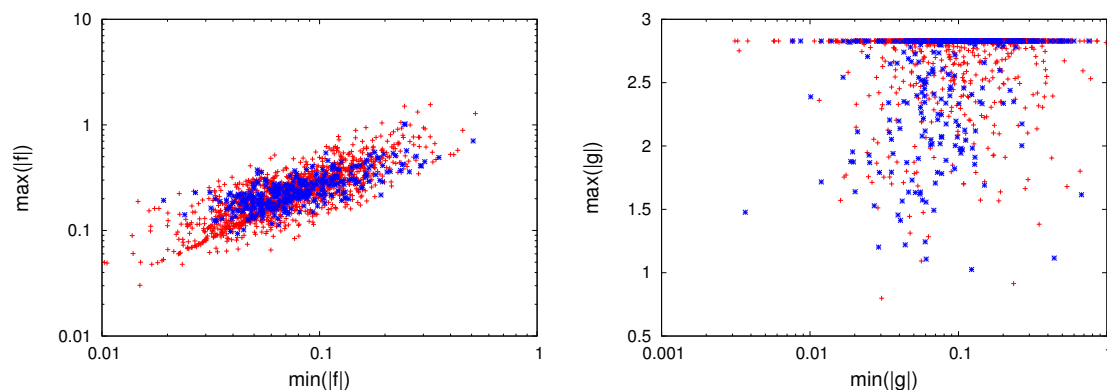


Figure 6. Viable regions of parameter space for the Yukawa couplings $f_{\alpha\beta}$ and $g_{i\alpha}$. Correct neutrino mass and mixing is obtained and flavor-changing constraints are satisfied. The blue benchmarks give the DM relic abundance in accordance with figure 4.

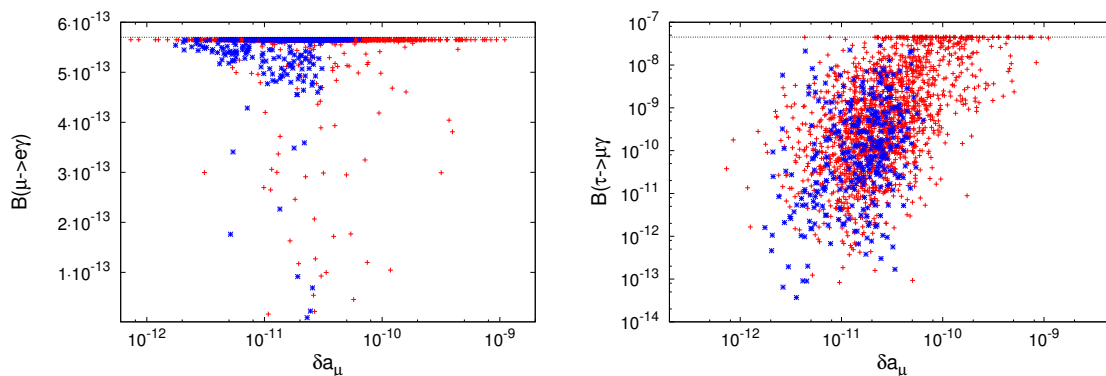


Figure 7. Branching fractions for lepton flavor violating decays versus the anomalous magnetic moment of the muon. The dashed lines represent the experimental upper bounds on the branching ratios.

In the above we employed the DM annihilation cross sections from section 5, finding $M_F \sim 6$ TeV. With this value of M_F , low-energy constraints are readily satisfied and viable neutrino masses are obtained. However, the (co-)annihilation cross sections are subject to a Sommerfeld enhancement due to $SU(2)_L$ gauge-boson exchange. This modifies the (co-)annihilation cross sections and increases M_F . When the enhancement is applied to s -wave (co-)annihilations via $SU(2)_L$ interactions, the requisite DM mass increases to $M_F \sim 10$ TeV [41]. In our case one should also include the enhancement for the p -wave annihilations [65], which is beyond the scope of this work. However, we anticipate similar results for our model and expect an $\mathcal{O}(1)$ correction to M_F due the enhancement. To determine if the model is likely to remain viable once the Sommerfeld effect is included, we studied the parameter space with heavier $M_F \lesssim 20$ TeV. We found that viable neutrino mass/mixings could be obtained while satisfying the various constraints. These results are already incorporated in the figures, as seen in figure 5, where $M_F \lesssim 20$ TeV is considered. These results indicate that the model should remain viable with $M_F \sim 10$ TeV.

7 Collider phenomenology

Although a detailed study of the collider phenomenology of our model is beyond the scope of the present work, we briefly discuss some important signatures at both the LHC and future e^+e^- colliders. If \mathcal{F}_1 provides the DM relic abundance one requires $M_f \sim 10$ TeV with $M_\phi > M_F$, placing both \mathcal{F} and ϕ well beyond the reach of foreseeable collider experiments. However, the singlet charged scalar S^\pm can remain within reach of TeV scale colliders. At the International Linear collider (ILC) [66], the charged scalars S^\pm can be directly produced through the t-channel process $e^+e^- \rightarrow S^+S^- \rightarrow l_\alpha^+ l_\beta^- + E_{miss}$, which includes lepton flavor violating final-states that can be observed as a pair of charged leptons with missing energy (similar to the KNT model [10]). However, due to different constraints in the KNT model, the corresponding charged scalar is not allowed to be as light as 300 GeV, as is the case here. Therefore it should be easier to test our model through this channel at the ILC for energies $\sqrt{s} = 500$ GeV and 1 TeV. At the LHC, this model can similarly be probed via the process $pp \rightarrow S^+S^- \rightarrow l_\alpha^+ l_\beta^- + E_{miss}$ with the charged scalars produced through Drell-Yan.

The region of parameter space with lighter values of $M_F \sim$ TeV is also of interest as it allows \mathcal{F} (and possibly ϕ) to be within reach of collider experiments like the LHC. In this parameter space \mathcal{F}_1 cannot provide the full DM relic abundance though it can provide a sub-leading contribution. The exotic fermions would be pair produced via weak interactions at the LHC as $pp \rightarrow W/Z \rightarrow \mathcal{F}\mathcal{F}$, with typical weak-scale cross sections (e.g. for $M_F \approx 300$ GeV, one expects a production cross section of $\mathcal{O}(10^2)$ fb at the 7 TeV LHC and $\mathcal{O}(10^3)$ fb at a 14 TeV LHC). Due to the exact (or approximate) Z_2 symmetry, the heavier fermions must decay weakly to lighter exotic fermions rather than directly to SM particles. For example, one could have the production process $pp \rightarrow W^+ \rightarrow \mathcal{F}^{++}\mathcal{F}^-$, with the charged fermions decaying via off-shell W bosons to leptonic final states as $\mathcal{F}^{++} \rightarrow \mathcal{F}^+ \ell^+ \nu_\ell$, and $\mathcal{F}^\pm \rightarrow \mathcal{F}^0 \ell^\pm \nu_\ell$, where $\ell = e, \mu, \tau$ denotes the SM lepton flavor. A typical final state would contain three charged leptons and missing energy, due to the DM and the neutrinos. Related final states with four charged leptons are also possible via $\mathcal{F}^{++}\mathcal{F}^{--}$ pair production.

8 Generalized KNT models

The model presented here is related to the proposal of KNT [6] and a recently discovered three-loop model with triplet fermions [11]. In this section we identify this relationship and show that the models form a larger set of generalized KNT models. Consider the loop diagram in figure 1. Adding S to the SM to allow the outer vertices, the choice for \mathcal{F} and ϕ is not unique. One can determine the basic conditions for a general fermion $\mathcal{F} \sim (1, R_{\mathcal{F}}, Y_{\mathcal{F}})$ and scalar $\phi \sim (1, R_{\phi}, Y_{\phi})$ that allow figure 1 to appear. The top vertex in figure 1 requires a term $\lambda_s(S^-)^2\phi^2 \subset V(H, \phi, S)$ in the potential. This fixes $Y_{\phi} = Y_S = 2$, which in turn fixes $Y_{\mathcal{F}} = -(Y_{\phi} + Y_{e_R}) = 0$. For even-valued $R_{\mathcal{F}}$, the model contains fractionally charged particles, the lightest of which is automatically stable and therefore excluded by cosmological constraints. Consequently only odd-valued $R_{\mathcal{F}}$ is viable, giving $R_{\mathcal{F}} = (2n + 1)$ for $n = 0, 1, \dots$. The $\overline{\mathcal{F}}\phi e_R$ vertex then fixes $R_{\phi} = R_{\mathcal{F}} = (2n + 1)$.

For $n = 0$ one has the KNT model, with $\mathcal{F} \sim (1, 1, 0)$ and $\phi \sim (1, 1, 2)$ [6], while $n = 1$ gives the recently-proposed triplet model with $\mathcal{F} \sim (1, 3, 0)$ and $\phi \sim (1, 3, 2)$ [11]. In both of these models one requires a new symmetry to remove the tree-level seesaw contributions. For $n = 2$ one obtains the present model, which gives $\mathcal{F} \sim (1, 5, 0)$ and $\phi \sim (1, 5, 2)$. Thus, $n = 2$ is the smallest value for which no symmetry is required to remove a tree-level seesaw mass — neutrino mass automatically appears at the three-loop level for $n \geq 2$, irrespective of whether a Z_2 symmetry is imposed.

We saw that DM longevity did not require a Z_2 symmetry in the $n = 2$ model due to the softly-broken accidental Z_2 symmetry (which becomes exact for $\lambda \rightarrow 0$). This feature is common for all even-valued n with $n \geq 2$, which is seen as follows. For all $n \geq 0$, the most-general Lagrangian seemingly contains the term $\lambda(S^-)\phi^* \times (\phi \times \phi)_{R_{\mathcal{F}}} \subset V(H, S, \phi)$, which breaks the Z_2 symmetry. Here $(\phi \times \phi)_{R_{\mathcal{F}}}$ denotes the SU(2)-contraction of $\phi \times \phi$ in the $R_{\mathcal{F}}$ representation; for odd-valued $R_{\mathcal{F}}$ this is always contained in the SU(2)-product, $R_{\mathcal{F}} \subset R_{\mathcal{F}} \times R_{\mathcal{F}}$. For $n < 2$, however, the models contain additional Z_2 symmetry breaking terms, including some that generate tree-level neutrino masses. On the other hand, for $n \geq 2$ the λ -term is the sole Z_2 symmetry breaking term. Thus, $n = 2$ marks the transition where the λ -term softly breaks the Z_2 symmetry, and all models with $n \geq 2$ seemingly possess a softly-broken accidental symmetry that becomes exact in the limit $\lambda \rightarrow 0$. However, although group theory gives $R_{\mathcal{F}} \times R_{\mathcal{F}} \supset R_{\mathcal{F}}$, the product $(\phi \times \phi)_{R_{\mathcal{F}}}$ in fact vanishes when the scalar is in the $R_{\mathcal{F}} = 2n + 1$ representation for odd-valued n .² Thus, for all even-valued $n \geq 2$, the models contain an accidental Z_2 symmetry that is softly broken by the term $\lambda(S^-)\phi^* \times (\phi \times \phi)_{R_{\mathcal{F}}} \subset V(H, S, \phi)$.

There is a very interesting by-product of these observations. If ϕ is in the $R_{\mathcal{F}} = 2n + 1$ representation with $n \geq 2$, the λ -term is the sole Z_2 symmetry breaking term in the model. However, for odd-valued n , the λ -term vanishes identically, and the accidental Z_2 symmetry becomes an exact symmetry of the full Lagrangian. Thus, for $R_{\mathcal{F}} = 7$, corresponding to $\mathcal{F} \sim (1, 7, 0)$ and $\phi \sim (1, 7, 2)$, one automatically obtains a model of radiative neutrino mass with a stable DM candidate due to an exact accidental Z_2 symmetry — no additional

²For distinct scalars ϕ and ϕ' , both in the $R_{\mathcal{F}} = 2n + 1$ representation, the SU(2) product $(\phi \times \phi)_{R_{\mathcal{F}}}$ is nonzero for all n . However, for identical scalars $\phi = \phi'$, one finds $(\phi \times \phi)_{R_{\mathcal{F}}} = 0$ for odd-valued n .

symmetry need be imposed.³ More generally, models with odd-valued $n > 2$ will generate neutrino mass and give stable DM candidates without invoking new symmetries.

9 Conclusion

We presented a three-loop model of neutrino mass whose most-general Lagrangian contains a softly-broken accidental Z_2 symmetry. In the limit that a single parameter vanishes, $\lambda \rightarrow 0$, the Z_2 symmetry becomes exact and the model contains a stable DM candidate. Even for nonzero $\lambda \ll 1$, however, the model can give a long-lived DM candidate. The model is related to the KNT model and its triplet variant, with the Z_2 symmetry being equivalent to the symmetry imposed in those models. In the present case, though, the symmetry is not needed to preclude tree-level neutrino mass, giving a viable model of neutrino mass irrespective of DM considerations. For sufficiently small λ , the model gives a unified solution to the DM and neutrino mass problems, with the novel feature of not requiring that a symmetry be imposed. We showed that neutrino mass can be generated and that important flavor-changing constraints can be satisfied. Taken purely as a neutrino mass model, the new physics can be $\mathcal{O}(\text{TeV})$, allowing the model to be explored at colliders. However, when DM is included the quintuplet fields must be heavy, with $M_F \sim 10 \text{ TeV}$, so that only the singlet scalar S can be within reach of colliders. None the less, the DM can be tested in future direct-detection experiments. We also noted interesting generalizations of this model in which DM stability results from an exact accidental symmetry, the simplest of which uses septuplet $SU(2)$ fields instead of quintuplets.

Acknowledgments

The authors thank T. Toma. AA is supported by the Algerian Ministry of Higher Education and Scientific Research under the CNEPRU Project No D01720130042. KM is supported by the Australian Research Council.

A Radiative neutrino mass

The Majorana neutrino masses are calculated to be

$$(\mathcal{M}_\nu)_{\alpha\beta} = \frac{5\lambda_S}{(4\pi^2)^3} \frac{m_\gamma m_\delta}{M_\phi} f_{\alpha\gamma} f_{\beta\delta} g_{\gamma i}^* g_{\delta i}^* \times F\left(\frac{M_i^2}{M_\phi^2}, \frac{M_S^2}{M_\phi^2}\right), \quad (\text{A.1})$$

where

$$F(\alpha, \beta) = \frac{\sqrt{\alpha}}{8\beta^2} \int_0^\infty dr \frac{r}{r+\alpha} \left(\int_0^1 dx \ln \frac{x(1-x)r + (1-x)\beta + x}{x(1-x)r + x} \right)^2. \quad (\text{A.2})$$

In obtaining this form of F we have neglected the lepton masses.

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³We thank T. Toma for communications on this point.

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