Inverse seesaw and dark matter in models with exotic lepton triplets

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ABSTRACT

We show that models with exotic leptons transforming as $E \sim (1, 3, -1)$ under the standard model gauge symmetry are well suited for generating neutrino mass via a radiative inverse seesaw. This approach realizes natural neutrino masses and allows multiple new states to appear at the TeV scale. The exotic leptons are therefore good candidates for new physics that can be probed at the LHC. Furthermore, remnant low-energy symmetries ensure a stable dark matter candidate, providing a link between dark matter and the origin of neutrino mass.

1. Introduction

The seesaw mechanism [1] provides a simple means for generating naturally light standard model (SM) neutrinos. This simplicity, however, comes at a cost. If the Dirac Yukawa couplings in the neutrino sector are large [e.g. $\mathcal{O}(1)$], as argued on the basis of naturalness, the gauge-singlet neutrinos must be very heavy, and there is essentially no hope of experimentally producing them. Thus, if the seesaw is realized naturally in nature, the origin of neutrino mass will remain a mystery for the foreseeable future.

Beyond the seesaw mechanism, however, there exist alternative approaches capable of explaining neutrino mass (see e.g. [2]). For example, light neutrinos can be realized naturally without invoking new high-energy physics if the right-handed neutrinos are low-energy composites of a confining hidden sector [3].[1] In principle, such models are easier to experimentally verify than a high-energy seesaw. In practice, models with light new physics can be similarly difficult to probe, due to very weak couplings. Indirect signals like flavor changing decays [5], CMB signals [6], or galactic X-ray signals [7], often provide the best way to probe such scenarios. This is more promising than the somewhat bleak outlook afforded by the canonical seesaw, though it would be preferable if the new physics associated with neutrino mass was amenable to direct experimental production.

The likelihood of discovering new-physics related to neutrino mass increases if it includes new TeV-scale particles that are charged under the SM. Radiative models of neutrino mass provide

a good example; here the SM couples to new fields that can potentially be probed at the LHC, while neutrino mass is generated as a small loop-effect [8]. These models are very promising from an experimental point of view, as the connection between small masses and new TeV-scale physics offers perhaps the best hope of discovery (see e.g. [9]).

With this in mind, it is worthwhile considering exotic states with $\mathcal{O}(\text{TeV})$ masses that have gauge invariant, renormalizable couplings to SM fermions. There are only three such fermionic extensions that can minimally couple to SM neutrinos via the Higgs [11]. The well-studied cases of the singlet neutrinos $\nu_R \sim (1, 1, 0)$, and the triplet fermions $\Sigma_R \sim (1, 3, 0)$, correspond to the type I [1] and type III [12] seesaws, respectively. Here we consider the relatively unexplored third case of the leptons $E_R = (E_R^0, E_R^-, E_R^{--})^T$, which form an $SU(2)_L$ triplet with hypercharge $Y = -1$. The phenomenology of these exotics has received some attention in the literature [10,11,13]; amongst other things, they induce new flavor-changing effects [11], and the doubly-charged fields have distinctive collider signatures [10,13].[2] Given that current bounds permit these exotics to be discovered at the LHC [10,11,13], it is worth asking if they enable new mechanisms for generating neutrino masses.

Perhaps surprisingly, not much attention has been given to this matter. Ref. [16] is the only prior work to consider the link between E_R and neutrino masses; they obtained masses of the form $m_\nu \simeq \lambda^3 \langle \phi^0 \rangle^4 / M_{\text{NP}}^3$, where $\langle \phi^0 \rangle$ is the SM scalar vacuum value, λ denotes generic dimensionless couplings, and M_{NP} is the scale of new physics. With this approach, one can only have $M_{\text{NP}} \simeq \mathcal{O}(\text{TeV})$

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[1] Equivalently, via AdS/CFT, if they propagate in a hidden warped space [4].

[2] The minimally coupled lepton doublet with $Y = -3/2$ also predicts a doubly charged lepton [10,14], but its phenomenology is markedly different. For related work see [15].

if the dimensionless couplings provide an additional tiny factor of $\mathcal{O}(10^{-10})$ [16].

In this work we consider alternative (radiative) models containing the fermions E_R. These models realize small neutrino masses without invoking tiny parameters, and permit the new physics to appear at the TeV scale. Furthermore, the models contain remnant low-energy symmetries that ensure the stability of a new dark matter candidate, providing a connection between neutrino mass and dark matter (for earlier works see [17]).

2. Neutrino mass and exotic triplets

We extend the SM to include the exotic triplets $E_{R,L}$, which transform under $G_{SM} = SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ as

$$E_{R,L} = \begin{pmatrix} E^0 \\ E^- \\ E^{--} \end{pmatrix}_{R,L} \sim (1,3,-1). \tag{1}$$

The gauge invariant Lagrangian thus includes the terms

$$\mathcal{L} \supset -Y_E \bar{L}_L \phi E_R - M_E \bar{E}_L E_R + \text{H.c.}, \tag{2}$$

where L_L (ϕ) is the SM lepton (Higgs) doublet. Symmetry breaking generates a mass between E_R^0 and the SM neutrinos. For a given value of M_E, the Yukawa couplings Y_E are constrained by Z decays and lepton flavor violation [11]. We are interested in $M_E \simeq 1$ TeV, in which case the matrix elements must satisfy [11]

$$(Y_E)_{ii} \lesssim 1 \quad \text{and} \quad (Y_E)_{ij} \lesssim 10^{-2} \quad \text{for } i \neq j, \tag{3}$$

barring tuned cancellations.

The extension by $E_{R,L}$ alone is insufficient to generate viable neutrino masses.[3] In Section 3 we will detail further extensions that permit viable masses, however, the resulting small neutrino masses can be understood prior to learning these details. In the models, the neutral-lepton mass matrix has the form

$$\mathcal{M}_{\text{inv}}^E = \begin{pmatrix} 0 & Y_E \langle \phi^0 \rangle & 0 \\ Y_E^\dagger \langle \phi^0 \rangle & \delta_R & M_E \\ 0 & M_E & \delta_L \end{pmatrix}, \tag{4}$$

in the basis $\{\nu_L, E_R^0, E_L^0\}$. Here, the entries $\delta_{L,R}$ are small and their (model-dependent) origin is explained in the next section, while the other entries follow from (2). This matrix (4) has the same form as that employed in the inverse seesaw mechanism [19]. Thus, writing $m = Y_E \langle \phi^0 \rangle$, the eigenvalues are of the standard inverse seesaw form:

$$m_1 = \frac{m^2 \delta_L}{m^2 + M_E^2} + \mathcal{O}(\delta^2),$$

$$m_{2,3} = \mp \sqrt{m^2 + M_E^2} + \frac{M_E^2 \delta_L}{2(m^2 + M_E^2)} + \frac{\delta_R}{2} + \mathcal{O}(\delta^2), \tag{5}$$

where $\delta_{L,R} \ll m, M_E$ is used. The light neutrino mass (i.e. m_1) is suppressed by the smallness of δ_L, with further suppression if $m \ll M_E$.

As with standard inverse-seesaw models, this scheme allows significant mixing between ν_L and the exotic (E_L), with $\theta \simeq \mathcal{O}(\sqrt{|m_1/\delta_L|})$. For $m_1 \lesssim 0.1$ eV and $\delta_L \lesssim 30$ keV (compatible with the new physics M_E being at the TeV scale), this mixing is

$\mathcal{O}(10^{-3})$ and non-unitary mixing effects may be detectable in future experiments [21]. Mixing between ν_L and E_R is at the level of $\mathcal{O}(m_1/m) \simeq \mathcal{O}(10^{-9})$ and remains negligible.[4]

Thus, with naturally small values of $\delta_{L,R}$, one can successfully obtain small neutrino masses when $E_{R,L}$ are at the TeV scale. Small values of $\delta_{L,R}$ can be obtained naturally if they arise as loop effects, resulting in a (non-standard) radiative inverse seesaw (as first considered in [20]). In the following section we present two models that realize this idea.

3. Radiative inverse seesaw with exotic lepton triplets

In what follows, we present two models that radiatively generate $\delta_{L,R} \neq 0$; one employing an extended gauge symmetry, the other a global $\mathbb{Z}_4$ symmetry.

3.1. Model with a gauged U(1)

Nonzero $\delta_{L,R}$ can be generated radiatively by introducing new particles that are charged under a hidden (or "dark") gauge symmetry. We consider the extended gauge group $G_{SM} \otimes U(1)_d$, with new fields that transform as

$$E_{R,L} \sim (1,3,-1)(0), \qquad N_{R,L} \sim (1,1,0)(d),$$

$$\xi \sim (1,3,-1)(-d), \qquad \eta \sim (1,1,0)(d), \qquad \chi \sim (1,1,0)(2d).$$

Here the first (second) line contains exotic fermions (scalars), and d is an arbitrary charge that can be set to unity. These particles, together with the SM fields (which are $U(1)_d$ singlets), form an anomaly-free set. In addition to (2), the Lagrangian contains the terms

$$\mathcal{L} \supset h_\xi \bar{E}_L \xi N_R + h_\chi \bar{N}_R (N_R)^c \chi + M_N \bar{N}_L N_R$$
$$+ h'_\xi \bar{E}_R \xi N_L + h'_\chi \bar{N}_L (N_L)^c \chi + \text{H.c.}, \tag{6}$$

where $h_{\xi,\chi}$ and $h'_{\xi,\chi}$ are Yukawa couplings.

Suppose that our model building choices enforce the following VEV pattern

$$\langle \phi^0 \rangle \simeq 174 \text{ GeV}, \qquad \langle \xi^0 \rangle, \langle \eta \rangle = 0, \qquad \langle \chi \rangle \neq 0. \tag{7}$$

Then, in the basis $\{\nu_L, E_R^0, E_L^0\}$, the tree-level mass matrix has the form (4) with a conserved lepton-number symmetry. In addition, the 2×2 mass matrix for the "dark" neutrinos $N_{R,L}$ decouples, and the associated mass eigenstates are two (Majorana) linear combinations of $N_{R,L}$. These dark neutrinos do not mix with the other fields – a statement that holds to all orders in perturbation theory, as we explain below.

The most general scalar potential is

$$V_S = (\phi^\dagger \phi)\big[\mu_\phi^2 + \lambda_\phi(\phi^\dagger \phi) + \lambda_{\phi\xi}(\xi^\dagger \xi) + \lambda_{\phi\eta}(\eta^\dagger \eta)$$
$$+ \lambda_{\phi\chi}(\chi^\dagger \chi)\big] + (\eta^\dagger \eta)\big[\mu_\eta^2 + \lambda_\eta(\eta^\dagger \eta) + \lambda_{\eta\chi}(\chi^\dagger \chi)\big]$$
$$+ (\xi^\dagger \xi)\big[\mu_\xi^2 + \lambda_\xi(\xi^\dagger \xi) + \lambda_{\xi\eta}(\eta^\dagger \eta) + \lambda_{\xi\chi}(\chi^\dagger \chi)\big]$$
$$+ (\chi^\dagger \chi)\big[\mu_\chi^2 + \lambda_\chi(\chi^\dagger \chi)\big]$$
$$+ \Big[\lambda_{\xi\phi\eta}\xi\phi\eta + \frac{1}{2}\mu_{\eta\chi}\eta\eta\chi^\dagger + \text{H.c.}\Big]. \tag{8}$$

Observe that lepton number symmetry is broken explicitly in the last line. As per usual in radiative models of neutrino mass, this

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[3] This is standard in radiative models, for which the extension by one exotic is typically insufficient to generate neutrino mass (see e.g. [8,18]).

[4] For these parameters, a normal type-I/III seesaw gives mixing of $\mathcal{O}(10^{-7})$. The type-II seesaw [22] does not lead to unitarity violation.
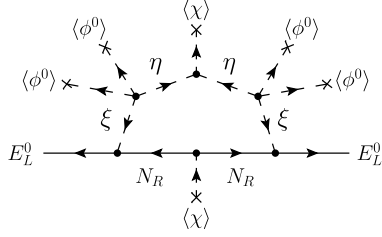
Fig. 1. One-loop diagram for the Majorana mass δ_L. Note that δ_R is obtained via M_E mass insertions at both ends.

explicit breaking of lepton number symmetry is radiatively communicated to the neutral lepton sector [8]. As a result, one expects the Majorana masses $\delta_{L,R}$ to be generated radiatively. Indeed, δ_L is generated via the one-loop dim-9 diagram in Fig. 1. The corresponding diagram for δ_R is obtained with appropriate mass insertions on the external lines.

Writing the scale of new physics as $M_{NP} \simeq M_\xi, M_\eta, M_N$ (the masses for $\xi, \eta, N_{R,L}$ respectively), the one-loop Majorana mass is approximately

$$\delta_L \simeq -\frac{1}{16\pi^2} \frac{h_\xi^2 h_\chi \lambda_{\xi\phi\eta}^2}{15} \frac{\langle\phi^0\rangle^4 \langle\chi\rangle^2}{M_{NP}^6} \mu_{\eta\chi}. \tag{9}$$

For $\langle\chi\rangle \simeq M_{NP} \simeq 1$ TeV, $\mu_{\eta\chi} \simeq 100$ GeV (electroweak scale) and $\lambda_{\xi\phi\eta}, h_{\xi,\chi} = \mathcal{O}(1)$, we have $|\delta_L| \simeq \mathcal{O}(10\text{--}100)$ keV. Using this in (5), with the conservative value $Y_E \simeq 10^{-2}$ for $M_E \simeq 1$ TeV, gives a light neutrino mass of about 10^{-1} eV. Hence, there is no severe tuning and small neutrino masses are compatible with the new physics appearing at the TeV scale.[5]

Note that the VEV pattern in (7) is easily achieved. A small hierarchy between $\langle\phi^0\rangle$ and $\langle\chi\rangle = \mathcal{O}(\text{TeV})$ is obtained provided the mixing term $\lambda_{\phi\chi}\phi^\dagger\phi\chi^\dagger\chi$ is not too large. The requirement that $\langle\xi^0\rangle, \langle\eta\rangle = 0$ is also compatible with $\langle\phi^0\rangle, \langle\chi\rangle \neq 0$, as long as all couplings involving ξ and η in V_S are positive. The stability of the classical vacua may be affected by radiative corrections, including fermion loops induced by $h_\xi, h'_\xi \neq 0$. However, for a given set of classical parameters that achieve (7), one can always select a suitable cut-off to ensure this minimum persists at the quantum level. If needed, one can also reduce h_ξ, h'_ξ to allow larger values of the cut-off. Note that smaller values of h_ξ help to suppress δ_L.

Since ξ is charged under G_{SM}, one might naively worry about the SM ρ-parameter. However, we require $\langle\xi^0\rangle = 0$, so there are no tree-level corrections to the SM value and the model lies safely within experimental limits. Note that symmetry breaking induces mass mixing between χ and ϕ^0, providing a channel for communication between the SM and the hidden sector. There is also mixing between ξ and η, though these scalars do not mix with χ and ϕ^0.

The hidden sector also communicates with the SM via gauge kinetic-mixing [23] between U(1)$_d$ and hypercharge. The mixing term in the Lagrangian is

$$\mathcal{L} \supset -\frac{\kappa}{2}(\cos\theta_w F_{\mu\nu} - \sin\theta_w Z_{\mu\nu})Z_d^{\mu\nu}, \tag{10}$$

where θ_w is the weak mixing angle, and the mixing parameter obeys $\kappa \ll 1$ (see below). To leading order in κ, the mixed kinetic-Lagrangian is diagonalized by the following field redefinitions:

$$A^\mu \to A^\mu - \kappa\cos\theta_w Z_d^\mu,$$

$$Z^\mu \to Z^\mu + \kappa\sin\theta_w Z_d^\mu,$$
$$Z_d^\mu \to Z_d^\mu. \tag{11}$$

In this basis, the mass terms for Z and Z_d are

$$\mathcal{L} \supset \frac{1}{2}m_Z^2 Z^2 + \kappa\sin\theta_w m_Z^2 Z Z_d + \frac{1}{2}m_{Z_d}^2 Z_d^2. \tag{12}$$

Hence, to this order, the Z–Z_d mixing angle is

$$\theta_d \simeq \kappa\sin\theta_w\left(\frac{m_Z^2}{m_Z^2 - m_{Z_d}^2}\right), \tag{13}$$

which is suppressed by both small κ and the mass ratio $m_Z^2/m_{Z_d}^2$. The corresponding mass eigenvalues are $M_Z^2 \simeq m_Z^2[1 + \mathcal{O}(\kappa^2)]$ and $M_{Z_d}^2 \simeq m_{Z_d}^2[1 + \mathcal{O}(\kappa^2)]$. With these results, one obtains the coupling between the *physical* massive vectors and matter:

$$\mathcal{L} \supset -Z^\mu[J_Z + \theta_d J_{Z_d}]_\mu - Z_d^\mu[J_{Z_d} - \kappa J_Y - \theta_d J_Z]_\mu,$$

where J_Y is the SM hypercharge current. Thus, the dominant effect of the kinetic mixing is to induce a small coupling of Z_d to J_Y when $\kappa \gg \theta_d$.

In addition to collider signals due to the triplets $E_{R,L}$, the new vector Z_d can produce observable effects. Detailed phenomenological studies of a hidden vector with mass $M_{Z_d} \gtrsim 200$ GeV have appeared in the literature [24]. In the present model, when Z_d and χ are the lightest exotic states, the constraints on κ are similar to those obtained in Ref. [24]:

$$|\kappa| \lesssim 0.006 \times \left(\frac{M_{Z_d}}{\text{TeV}}\right). \tag{14}$$

Clearly, $\kappa \ll 1$ is necessary for $M_{Z_d} \lesssim$ TeV, though the constraint is compatible with κ being generated at the one-loop level.

Once produced in a collider, the hidden vector can decay back to SM fermions. These decays proceed through the kinetic mixing, and in the limit $M_{Z_d} \gg M_Z$ the widths are

$$\Gamma(Z_d \to \bar{f}f) \simeq \frac{\kappa^2\alpha N_c}{6\cos^2\theta_w}[Y_{f_L}^2 + Y_{f_R}^2]M_{Z_d}. \tag{15}$$

Here α is the fine-structure constant, Y_f is the hypercharge value, and N_c is the color factor. Notice the distinctive hypercharge dependence, which ensures that the branching fractions to SM fermions are approximately set by $[Y_{f_L}^2 + Y_{f_R}^2]$. Summing over kinematically available final states (including W^+W^-) gives [24]

$$\Gamma_{Z_d} \approx 0.2 \times \left(\frac{\kappa^2}{0.01}\right)\left(\frac{M_{Z_d}}{\text{TeV}}\right) \text{GeV}. \tag{16}$$

The $\mathcal{O}(\kappa^2)$ suppression forces the width to be rather narrow, but discovery remains possible in regions of parameter space [24].

Finally, let us point out another welcome feature of the model — namely the existence of a stable dark matter candidate. Note that when χ develops a nonzero VEV, the U(1)$_d$ symmetry is broken to a residual $\mathbb{Z}_2$. Since $\langle\xi^0\rangle, \langle\eta\rangle = 0$, this symmetry, whose action is defined by

$$\{N_{R,L}, \xi, \eta\} \to -\{N_{R,L}, \xi, \eta\}, \qquad \psi_{\text{others}} \to \psi_{\text{others}}, \tag{17}$$

remains as an exact symmetry of the low-energy Lagrangian. This prevents $N_{R,L}$ from mixing with the other leptons, as aforementioned. Furthermore, it forces the lightest particle among $N_{R,L}, \xi^0$ and η to be absolutely stable, ensuring a dark matter candidate. Since neutrino masses can be generated with couplings as large as $\mathcal{O}(1)$, it is trivial to obtain the correct dark matter abundance via freeze-out with $M_{NP} \simeq \mathcal{O}(\text{TeV})$. The dark matter is therefore a

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[5] The model also generates a three-loop dim-7 diagram for the $\bar{\nu}_L(\nu_L)^c$ mass term. We show in Appendix A that this does not affect our conclusions. There is also a two-loop dim-7 diagram for $\delta_{L,R}$, but this is subdominant when the new states are $\mathcal{O}(\text{TeV})$ (see Section 3.2).

standard WIMP whose stability is related to the origin of neutrino mass (see e.g. [17]).

The charged components of ξ also transform non-trivially in (17). Thus, one must ensure that ξ^- and ξ^{--} can decay in order to avoid having a (cosmologically excluded) stable charged particle. This is easy to achieve. For example, if the dark matter candidate is the lightest linear-combination of $N_{R,L}$, the decay $\xi^- \to E^- N$ is allowed provided $M_\xi > M_E + M_N$. Similarly, if χ is the dark matter candidate, decays like $\xi^- \to \chi W^- \gamma$ can be allowed.[6]

Note that the phenomenology of Z_d also depends on the mass spectrum of the exotics. For instance, when the dark matter candidate is the lightest linear-combination of $N_{R,L}$, the invisible decay $Z_d \to NN$ may occur. The width for this decay is $\Gamma(Z_d \to NN) \sim g_d^2 M_{Z_d}$, which is expected to dominate the SM width in (15), given the $\mathcal{O}(\kappa^2)$ suppression of the latter. Consequently, Z_d decays produce missing energy signals in this case. Also note that N would be an inelastic dark matter candidate, so that direct-detection bounds from XENON100 [25] would not be a concern.

3.2. Model with a global $\mathbb{Z}_4$

We have seen that a simple gauge extension of the SM allows one to generate SM neutrino masses via the exotic leptons $E_{R,L}$. In this subsection we briefly show that a related approach can be employed in models with a discrete symmetry. Specifically, we consider a global $\mathbb{Z}_4$ symmetry, with exotic fields that transform under $G_{SM} \times \mathbb{Z}_4$ as

$$E_{R,L} \sim (1,3,-1)(1), \qquad N_R \sim (1,1,0)(i),$$
$$\xi \sim (1,3,-1)(-i), \qquad \eta \sim (1,1,0)(i), \qquad \chi \sim (1,1,0)(-1). \tag{18}$$

As before, $E_{R,L}$ and N_R are fermions, while ξ, η and χ are scalars. All SM particles are neutral under $\mathbb{Z}_4$. Note that the chiral state N_L is no longer required, as the leptons in (18), together with the SM fermions, form an anomaly-free set. As a result, the interaction Lagrangian does not contain the N_L-terms shown in (6) but is otherwise identical. The scalar potential remains as in (8), with lepton number symmetry broken explicitly, so one expects nonzero $\delta_{L,R}$ to be generated radiatively.

In addition to the change in particle content, we consider the following VEV pattern

$$\langle \phi^0 \rangle \simeq 174 \text{ GeV}, \qquad \langle \xi^0 \rangle, \langle \eta \rangle, \langle \chi \rangle = 0. \tag{19}$$

Observe that the $\mathbb{Z}_4$-charged scalars have vanishing VEVs so the discrete symmetry remains intact. This prevents any mixing between N_R and the other leptons. It also forbids the one-loop diagram for δ_L in Fig. 1. However, nonzero δ_L is generated at the two-loop level by the diagram in Fig. 2.

In the simplifying limit $M_\chi \gg M_{NP}$, where M_{NP} is the (approximate) common mass of the other exotics, one gets[7]

$$\delta_L \simeq -\frac{1}{(16\pi^2)^2} \left(h_\xi^2 h_\chi \lambda_{\xi\phi\eta}^2 \right) \frac{\langle \phi^0 \rangle^4}{M_\chi^2 M_{NP}^2} \mu_{\chi\eta}. \tag{20}$$

Thus, one readily obtains suitable light neutrino masses for natural parameters.

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[6] Treating the ξ^0–η mixing as a mass insertion, the components of ξ are degenerate at the classical level. Loop corrections are expected to lift this degeneracy. If the radiative splittings make the charged components heavier, decays like $\xi^- \to \xi^0 W^- \gamma$ would also occur (the W could be off-shell).

[7] With Feynman parameters one can relate the two-loop integral to that evaluated in [26].
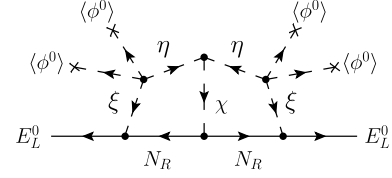

Fig. 2. Two-loop diagram for the Majorana mass δ_L. δ_R is obtained via M_E mass insertions at both ends.

Unsurprisingly, the exact $\mathbb{Z}_4$ symmetry ensures that the lightest $\mathbb{Z}_4$-charged state is a stable dark matter candidate. Distinct from the U(1)$_d$ model, however, one can now have a multi-component dark matter scenario. For example, if χ is the lightest $\mathbb{Z}_4$-charged state, then the lightest field out of η, ξ and N_R is also stable. Similarly, if either of η, ξ or N_R are the lightest $\mathbb{Z}_4$-charged state, then χ is also stable if $M_\chi < 2M_{DM}$ (M_{DM} being the mass of the lightest dark matter candidate). The phenomenology of this model differs from the U(1)$_d$ model due to the absence of Z_d (and N_L), but the model is otherwise very similar.

4. Conclusion

We have presented simple models for neutrino mass that realize a radiative inverse seesaw via the exotic lepton triplets $E \sim (1,3,-1)$. The models employ an extended field content, and allow multiple new states to appear at the TeV scale. The leptons E, and the hidden vector Z_d (appearing in one model), are good candidates for new physics that can probed at the LHC. Furthermore, the models contain stable dark matter candidates due to remnant low-energy symmetries.

Appendix A

Adding a nonzero $\bar{\nu}_L (\nu_L)^c$ mass term to (4) gives

$$\mathcal{M}_{inv} = \begin{pmatrix} \delta_m & m & 0 \\ m & \delta_R & M_E \\ 0 & M_E & \delta_L \end{pmatrix}. \tag{A.1}$$
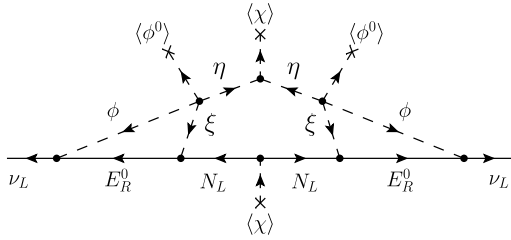
In the small δ_m limit, the light eigenvalue is now given by

$$m_1 = \underbrace{\frac{\delta_m M_E^2}{m^2 + M_E^2}}_{\simeq \delta_m} + \frac{\delta_L m^2}{m^2 + M_E^2} + \mathcal{O}(\delta^2),$$

while the heavier ones remain as $m_{2,3} = \mp\sqrt{m^2 + M_E^2} + \mathcal{O}(\delta)$. Thus, provided $\delta_m \gg m^2 \delta_L/(m^2 + M_E^2)$, the lightest eigenvalue remains on the order of the standard inverse seesaw result, $m_1 = \mathcal{O}(m^2 \delta_L/M_E^2)$.

In the U(1)$_d$ model of Section 3.1, nonzero δ_m is generated at the three-loop level (see Fig. 3). In the large M_ξ limit one can approximate this diagram as

$$\delta_m \simeq \frac{8h_\xi^2 h_\chi \lambda_{\xi\phi\eta}^2}{(16\pi^2)^3} \frac{Y_E^2 \langle \phi^0 \rangle^2 \langle \chi \rangle^2}{3M_\xi^4} \mu_{\chi\eta}$$
$$\times \left\{ \log\left[\frac{M_\xi^2 + M_{NP}^2}{M_{NP}^2} \right] - 1 \right\}^2, \tag{A.2}$$

Fig. 3. A three-loop dim-7 contribution to the $\bar{\nu}_L(\nu_L)^c$ entry of the neutral lepton mass matrix in the U(1)$_d$ model.

where M_{NP} denotes the mass of the other exotics. For $M_{\mathrm{NP}} \lesssim$ TeV, the result (A.2) can be subdominant to, or on the same order as, the one-loop expression obtained in Section 3.1, due to the additional loop-factor suppression.

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