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Scotogenic inverse seesaw model of neutrino mass

Sean Fraser, Ernest Ma *, Oleg Popov

Department of Physics and Astronomy, University of California, Riverside, CA 92521, USA

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ABSTRACT

A variation of the original 2006 radiative seesaw model of neutrino mass through dark matter is shown to realize the notion of inverse seesaw naturally. The dark-matter candidate here is the lightest of three real singlet scalars which may also carry flavor.

In 1998, the simplest realizations of the dimension-five operator [1] for Majorana neutrino mass, i.e. $(\nu_i\phi^0)(\nu_j\phi^0)$, were discussed systematically [2] for the first time. Not only was the nomenclature for the three and only three tree-level seesaw mechanisms established: (I) heavy singlet neutral Majorana fermion N [3], (II) heavy triplet Higgs scalar (ξ^{++}, ξ^+, ξ^0) [4], and (III) heavy triplet Majorana fermion $(\Sigma^+, \Sigma^0, \Sigma^-)$ [5], the three generic one-loop irreducible radiative mechanisms involving fermions and scalars were also written down for the first time. Whereas one such radiative mechanism was already well-known since 1980, i.e. the Zee model [6], a second was not popularized until eight years later in 2006, when it was used [7] to link neutrino mass with dark matter, called *scotogenic* from the Greek *scotos* meaning darkness. The third remaining unused mechanism is the subject of this paper. It will be shown how it is a natural framework for a scotogenic inverse seesaw model of neutrino mass, as shown in Fig. 1. The new particles are three real singlet scalars $s_{1,2,3}$, and one set of doublet fermions $(E^0, E^-)_{L,R}$, and one Majorana singlet fermion N_L, all of which are odd under an exactly conserved discrete symmetry Z_2. This specific realization was designated T1-3-A with $\alpha = 0$ in the compilation of Ref. [8]. Note however that whereas $(E^0, E^-)_L$ is not needed to complete the loop, it serves the dual purpose of (1) rendering the theory to be anomaly-free and (2) allowing E to have an invariant mass for the implementation of the inverse seesaw mechanism.

The notion of inverse seesaw [9–11] is based on an extension of the 2×2 mass matrix of the canonical seesaw to a 3×3 mass matrix by the addition of a second singlet fermion. In the space spanned by (ν, N, S), where ν is part of the usual lepton doublet (ν, l) and N, S are singlets, all of which are considered left-handed, the most general 3×3 mass matrix is given by

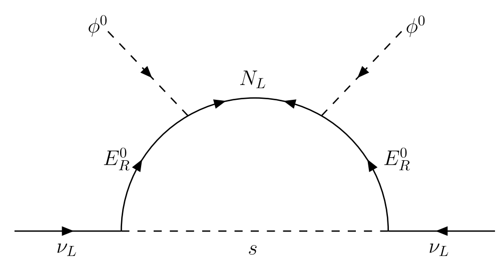

Fig. 1. One-loop generation of inverse seesaw neutrino mass.

$$\mathcal{M}_\nu = \begin{pmatrix} 0 & m_2 & 0 \\ m_2 & m_N & m_1 \\ 0 & m_1 & m_S \end{pmatrix}. \tag{1}$$

The zero $\nu - S$ entry is justified because there is only one ν to which N and S may couple through the one Higgs field ϕ^0. The linear combination which couples may then be redefined as N, and the orthogonal combination which does not couple is S. If $m_{S,N}$ is assumed much less than m_1, then the induced neutrino mass is

$$m_\nu \simeq \frac{m_2^2 m_S}{m_1^2}. \tag{2}$$

This formula shows that a nonzero m_ν depends on a nonzero m_S, and a small m_ν is obtained by a combination of small m_S and m_2/m_1. This is supported by the consideration of an approximate symmetry, i.e. lepton number L, under which $\nu, S \sim +1$ and $N \sim -1$. Thus $m_{1,2}$ conserve L, but m_S breaks it softly by 2 units. Note that there is also a finite one-loop contribution from m_N [12,13].

Other assumptions about m_1, m_S, m_N are also possible [14]. If $m_2, m_N \ll m_1^2/m_S$ and $m_1 \ll m_S$, then a double seesaw occurs with the same formula as that of the inverse seesaw, but of course with

* Corresponding author.

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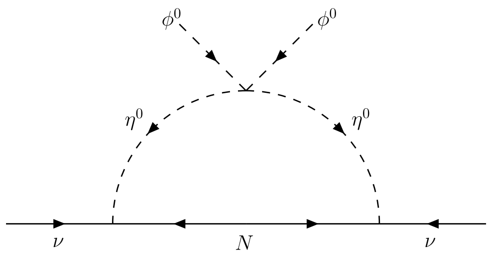

Fig. 2. One-loop generation of seesaw neutrino mass with heavy Majorana N.

a different mass hierarchy. If $m_1, m_2 \ll m_N$ and $m_1^2/m_N \ll m_S \ll m_1$, then a lopsided seesaw [14] occurs with $m_\nu \simeq -m_2^2/m_N$ as in the canonical seesaw, but $\nu - S$ mixing may be significant, i.e. $m_1 m_2/m_S m_N$, whereas $\nu - N$ mixing is the same as in the canonical seesaw, i.e. $\sqrt{m_\nu/m_N}$. In the inverse seesaw, $\nu - N$ mixing is even smaller, i.e. m_ν/m_2, but $\nu - S$ mixing is much larger, i.e. m_2/m_1, which is only bounded at present by about 0.03 [15]. In the double seesaw, the effective mass of N is m_1^2/m_S, so $\nu - N$ mixing is also $\sqrt{m_\nu/m_N}$. Here $m_S \gg m_N$, so the $\nu - S$ mixing is further suppressed by m_1/m_S.

In the original scotogenic model [7], neutrino mass is radiatively induced by heavy neutral Majorana singlet fermions $N_{1,2,3}$ as shown in Fig. 2. However, they may be replaced by Dirac fermions. In that case, a $U(1)_D$ symmetry may be defined [16], under which $\eta_{1,2}$ transform oppositely. If Z_2 symmetry is retained, then a radiative inverse seesaw neutrino mass is also possible [17,18]. We discuss here instead the new mechanism of Fig. 1, based on the third one-loop realization of neutrino mass first presented in Ref. [2]. The smallness of m_N, i.e. the Majorana mass of N_L, may be naturally connected to the violation of lepton number by two units, as in the original inverse seesaw proposal using Eq. (1). It may also be a two-loop effect as first proposed in Ref. [19], with a number of subsequent papers by other authors, including Refs. [20–22].

In our model, lepton number is carried by $(E^0, E^-)_{L,R}$ as well as N_L. This means that the Yukawa term $\bar{N}_L(E_R^0\phi^0 - E_R^-\phi^+)$ is allowed, but not $N_L(E_L^0\phi^0 - E_L^-\phi^+)$. In the 3×3 mass matrix spanning $(\bar{E}_R^0, E_L^0, N_L)$, i.e.

$$\mathcal{M}_{E,N} = \begin{pmatrix} 0 & m_E & m_D \\ m_E & 0 & 0 \\ m_D & 0 & m_N \end{pmatrix}, \tag{3}$$

m_E comes from the invariant mass term $(\bar{E}_R^0 E_L^0 + E_R^+ E_L^-)$, m_D comes from the Yukawa term given above connecting N_L with E_R^0 through $\langle\phi^0\rangle = v$, and m_N is the soft lepton-number breaking Majorana mass of N_L. Assuming that $m_N \ll m_D, m_E$, the mass eigenvalues of $\mathcal{M}_{E,N}$ are

$$m_1 = \frac{m_E^2 m_N}{m_E^2 + m_D^2}, \tag{4}$$

$$m_2 = \sqrt{m_E^2 + m_D^2} + \frac{m_D^2 m_N}{2(m_E^2 + m_D^2)}, \tag{5}$$

$$m_3 = -\sqrt{m_E^2 + m_D^2} + \frac{m_D^2 m_N}{2(m_E^2 + m_D^2)}. \tag{6}$$

In the limit $m_N \to 0$, E_R^0 pairs up with $E_L^0 \cos\theta + N_L \sin\theta$ to form a Dirac fermion of mass $\sqrt{m_E^2 + m_D^2}$, where $\sin\theta = m_D/\sqrt{m_E^2 + m_D^2}$. This means that the one-loop integral of Fig. 1 is well approximated by

$$m_\nu = \frac{f^2 m_D^2 m_N}{16\pi^2(m_E^2 + m_D^2 - m_s^2)}\left[1 - \frac{m_s^2 \ln((m_E^2 + m_D^2)/m_s^2)}{(m_E^2 + m_D^2 - m_s^2)}\right]. \tag{7}$$

This expression is indeed of the form expected of the inverse seesaw.

The radiative mechanism of Fig. 1 is also suitable for supporting a discrete flavor symmetry, such as Z_3. Consider the choice

$$\begin{aligned} &(\nu_i, l_i)_L \sim \underline{1}, \underline{1}', \underline{1}'', \qquad s_1 \sim \underline{1}, \\ &(s_2 + i s_3)/\sqrt{2} \sim \underline{1}', \qquad (s_2 - i s_3)/\sqrt{2} \sim \underline{1}'', \end{aligned} \tag{8}$$

with mass terms $m_s^2 s_1^2 + m_s'^2(s_2^2 + s_3^2)$, then the induced 3×3 neutrino mass matrix is of the form

$$\begin{aligned} \mathcal{M}_\nu &= \begin{pmatrix} f_e & 0 & 0 \\ 0 & f_\mu & 0 \\ 0 & 0 & f_\tau \end{pmatrix} \begin{pmatrix} I(m_s^2) & 0 & 0 \\ 0 & 0 & I(m_s'^2) \\ 0 & I(m_s'^2) & 0 \end{pmatrix} \\ &\quad \times \begin{pmatrix} f_e & 0 & 0 \\ 0 & f_\mu & 0 \\ 0 & 0 & f_\tau \end{pmatrix} \\ &= \begin{pmatrix} f_e^2 I(m_s^2) & 0 & 0 \\ 0 & 0 & f_\mu f_\tau I(m_s'^2) \\ 0 & f_\mu f_\tau I(m_s'^2) & 0 \end{pmatrix}, \end{aligned} \tag{9}$$

where I is given by Eq. (7) with f^2 removed. Let $l_{iR} \sim \underline{1}, \underline{1}', \underline{1}''$, then the charged-lepton mass matrix is diagonal using just the one Higgs doublet of the standard model, in keeping with the recent discovery [23,24] of the 125 GeV particle. To obtain a realistic neutrino mass matrix, we break Z_3 softly, i.e. with an arbitrary 3×3 mass-squared matrix spanning $s_{1,2,3}$, which leads to

$$\begin{aligned} &\begin{pmatrix} 1 & 0 & 0 \\ 0 & 1/\sqrt{2} & i/\sqrt{2} \\ 0 & 1/\sqrt{2} & -i/\sqrt{2} \end{pmatrix} O^T \begin{pmatrix} I(m_{s1}^2) & 0 & 0 \\ 0 & I(m_{s2}^2) & 0 \\ 0 & 0 & I(m_{s3}^2) \end{pmatrix} \\ &\quad \times O \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1/\sqrt{2} & 1/\sqrt{2} \\ 0 & i/\sqrt{2} & -i/\sqrt{2} \end{pmatrix}, \end{aligned} \tag{10}$$

where O is an orthogonal matrix but not the identity, and there can be three different mass eigenvalues $m_{s1,s2,s3}$ for the $s_{1,2,3}$ sector. The assumption of Eq. (8) results in Eq. (10) and allows the following interesting pattern for the neutrino mass matrix $\mathcal{M}_\nu$. The Yukawa couplings $f_{e,\mu,\tau}$ may be rendered real by absorbing their phases into the arbitrary relative phases between E_R^0 and $\nu_{e,\mu,\tau}$. If we further assume $f_\mu = f_\tau$, then $\mathcal{M}_\nu$ is of the form [25]

$$\mathcal{M}_\nu = \begin{pmatrix} A & C & C^* \\ C & D^* & B \\ C^* & B & D \end{pmatrix}, \tag{11}$$

where A and B are real. Note that this pattern is protected by a symmetry first pointed out in Ref. [26], i.e. $e \to e$ and $\mu - \tau$ exchange with CP conjugation, and appeared previously in Refs. [27,28]. As such, it is also guaranteed to yield maximal $\nu_\mu - \nu_\tau$ mixing ($\theta_{23} = \pi/4$) and maximal CP violation, i.e. $\exp(-i\delta) = \pm i$, whereas θ_{13} may be nonzero and arbitrary. Our scheme is thus a natural framework for this possibility. Further, from Eq. (7), it is clear that it is also a natural framework for quasi-degenerate neutrino masses as well. Let

$$F(x) = \frac{1}{1-x}\left[1 + \frac{x\ln x}{1-x}\right], \tag{12}$$

where $x = m_s^2/(m_E^2 + m_D^2)$, then Eq. (7) becomes

$$m_\nu = \frac{f^2 m_D^2 m_N}{(m_E^2 + m_D^2)} F(x). \tag{13}$$

Since $F(0) = 1$ and goes to zero only as $x \to \infty$, this scenario does not favor a massless neutrino. If $f_{e,\mu,\tau}$ are all comparable in magnitude, the most likely outcome is three massive neutrinos with comparable masses.

Since the charged leptons also couple to $s_{1,2,3}$ through E^-, there is an unavoidable contribution to the muon anomalous magnetic moment given by [29]

$$\Delta a_\mu = \frac{(g-2)_\mu}{2} = \frac{f_\mu^2 m_\mu^2}{16\pi^2 m_E^2} \sum_i |U_{\mu i}|^2 G(x_i), \tag{14}$$

where

$$G(x) = \frac{1 - 6x + 3x^2 + 2x^3 - 6x^2 \ln x}{6(1-x)^4}, \tag{15}$$

with $x_i = m_{si}^2/m_E^2$ and

$$U = O \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1/\sqrt{2} & 1/\sqrt{2} \\ 0 & i/\sqrt{2} & -i/\sqrt{2} \end{pmatrix}. \tag{16}$$

To get an estimate of this contribution, let $x_i \ll 1$, then $\Delta a_\mu = f_\mu^2 m_\mu^2/96\pi^2 m_E^2$. For $m_E \sim 1$ TeV, this is of order $10^{-11} f_\mu^2$, which is far below the present experimental sensitivity of 10^{-9} and can be safely ignored. The related amplitude for $\mu \to e\gamma$ is given by

$$A_{\mu e} = \frac{e f_\mu f_e m_\mu}{32\pi^2 m_E^2} \sum_i U_{ei}^* U_{\mu i} G(x_i). \tag{17}$$

Using the most recent $\mu \to e\gamma$ bound [30]

$$B = \frac{12\pi^2 |A_{\mu e}|^2}{m_\mu^2 G_F^2} < 5.7 \times 10^{-13}, \tag{18}$$

and the approximation $\sum_i U_{ei}^* U_{\mu i} G(x_i) \sim 1/36$ (based on tribimaximal mixing with $x_1 \sim 0$ and $x_2 \sim 1$) and $m_E \sim 1$ TeV, we find

$$f_\mu f_e < 0.03. \tag{19}$$

Let $f_{e,\mu,\tau} \sim 0.1$, $m_N \sim 10$ MeV, $m_D \sim 10$ GeV, $m_E \sim 1$ TeV, then the very reasonable scale of $m_\nu \sim 0.1$ eV in Eq. (7) is obtained, justifying its inverse seesaw origin. Since N_L is the lightest particle with odd Z_2, it is a would-be dark matter candidate. However, suppose we add N_R so that the two pair up to have a large invariant Dirac mass, then the lightest scalar (call it S) among $s_{1,2,3}$ is a dark-matter candidate. It interacts with the standard-model Higgs boson h according to

$$-\mathcal{L}_{int} = \frac{\lambda_{hS}}{2} v h S^2 + \frac{\lambda_{hS}}{4} h^2 S^2. \tag{20}$$

If we assume that all its other interactions are suppressed, then the annihilations $SS \to h \to$ SM particles and $SS \to hh$ determine its relic abundance, whereas its elastic scattering off nuclei via h exchange determines its possible direct detection in underground experiments. A detailed analysis [31] shows that the present limit of the invisible width of the observed 125 GeV particle (identified as h) allows m_S to be only within several GeV below $m_h/2$ or greater than about 150 GeV using the recent LUX data [32]. Note that the vector fermion doublet (E^0, E^-) is not the usually considered vector lepton doublet because it is odd under Z_2 and cannot mix with the known leptons.

In conclusion, we have shown how neutrino mass and dark matter may be connected using a one-loop mechanism proposed in 1998. This scotogenic model is naturally suited to implement the notion of inverse seesaw for neutrino mass, allowing the scale of new physics to be 1 TeV or less. The imposition of a softly broken Z_3 flavor symmetry yields an interesting pattern of radiative neutrino mass, allowing for maximal θ_{23} and maximal *CP* violation. The real singlet scalars in the dark sector carry lepton flavor, the lightest of which is absolutely stable. Our proposal provides thus a natural theoretical framework for this well-studied phenomenological possibility.

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References

[1] S. Weinberg, Phys. Rev. Lett. 43 (1979) 1566.
[2] E. Ma, Phys. Rev. Lett. 81 (1998) 1171.
[3] T. Yanagida, in: O. Sawada, A. Sugamoto (Eds.), Proc. of the Workshop on Unified Theories and Baryon Number in the Universe, KEK, Tsukuba, 1979, p. 95;
M. Gell-Mann, P. Ramond, R. Slansky, in: P. van Nieuwenhuizen, D.Z. Freedman (Eds.), Supergravity, North-Holland, Amsterdam, 1979, p. 315;
R.N. Mohapatra, G. Senjanovic, Phys. Rev. Lett. 44 (1980) 912;
P. Minkowski, Phys. Lett. B 67 (1977) 421.
[4] J. Schechter, J.W.F. Valle, Phys. Rev. D 22 (1980) 2227;
M. Magg, C. Wetterich, Phys. Lett. B 94 (1980) 61;
T.P. Cheng, L.-F. Li, Phys. Rev. D 22 (1980) 2860;
R.N. Mohapatra, G. Senjanovic, Phys. Rev. D 23 (1981) 165.
[5] R. Foot, H. Lew, X.-G. He, G.C. Joshi, Z. Phys. C 44 (1989) 441.
[6] A. Zee, Phys. Lett. B 93 (1980) 389.
[7] E. Ma, Phys. Rev. D 73 (2006) 077301.
[8] D. Restrepo, O. Zapata, C.E. Yaguna, J. High Energy Phys. 1311 (2013) 011.
[9] D. Wyler, L. Wolfenstein, Nucl. Phys. B 218 (1983) 205.
[10] R.N. Mohapatra, J.W.F. Valle, Phys. Rev. D 34 (1986) 1642.
[11] E. Ma, Phys. Lett. B 191 (1987) 287.
[12] D. Aristizabal Sierra, C.E. Yaguna, J. High Energy Phys. 1108 (2011) 013.
[13] P.S. Bhupal Dev, A. Pilaftsis, Phys. Rev. D 86 (2012) 113001.
[14] E. Ma, Mod. Phys. Lett. A 24 (2009) 2161.
[15] A. de Gouvea, et al., arXiv:1310.4340 [hep-ex].
[16] E. Ma, I. Picek, B. Radovcic, Phys. Lett. B 726 (2013) 744.
[17] T. Hambye, K. Kannike, E. Ma, M. Raidal, Phys. Rev. D 75 (2007) 095003.
[18] H. Okada, T. Toma, Phys. Rev. D 86 (2012) 033011.
[19] E. Ma, Phys. Rev. D 80 (2009) 013013.
[20] F. Bazzocchi, Phys. Rev. D 83 (2011) 093009.
[21] S.S.C. Law, K.L. McDonald, Phys. Lett. B 713 (2012) 490.
[22] G. Guo, X.-G. He, G.-N. Li, J. High Energy Phys. 1210 (2012) 044.
[23] ATLAS Collaboration, G. Aad, et al., Phys. Lett. B 716 (2012) 1.
[24] CMS Collaboration, S. Chatrchyan, et al., Phys. Lett. B 716 (2012) 30.
[25] E. Ma, Phys. Rev. Lett. 112 (2014) 091801.
[26] W. Grimus, L. Lavoura, Phys. Lett. B 579 (2004) 113.
[27] E. Ma, Phys. Rev. D 66 (2002) 117301.
[28] K.S. Babu, E. Ma, J.W.F. Valle, Phys. Lett. B 552 (2003) 207.
[29] S. Kanemitsu, K. Tobe, Phys. Rev. D 86 (2012) 095025.
[30] MEG Collaboration, J. Adams, et al., Phys. Rev. Lett. 110 (2013) 201801.
[31] J.M. Cline, P. Scott, K. Kainulainen, C. Weniger, Phys. Rev. D 88 (2013) 055025.
[32] LUX Collaboration, D.S. Akerib, et al., Phys. Rev. Lett. 112 (2014) 091303.