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Radiative Scaling Neutrino Mass with A_4 Symmetry

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Abstract

A new idea for neutrino mass was proposed recently, where its smallness is not due to the seesaw mechanism, i.e. not inversely proportional to some large mass scale. It comes from a one-loop mechanism with dark matter in the loop consisting of singlet Majorana fermions N_i with masses of order 10 keV and neutrino masses are scaled down from them by factors of about 10^{-5}. We discuss how this model may be implemented with the non-Abelian discrete symmetry A_4 for neutrino mixing, and consider the phenomenology of N_i as well as the extra scalar doublet (η^+, η^0).

The origin of neutrino mass is the topic of many theoretical discussions. The consensus is that its smallness is due to some mass scale larger than the electroweak breaking scale of about 100 GeV. If there are no particles beyond those of the standard model (SM) lighter than this scale, then the well-known unique dimension-five opeartor [1]

$$\mathcal{L}_5 = \frac{-f_{ij}}{2\Lambda}(\nu_i \phi^0 - l_i \phi^+)(\nu_j \phi^0 - l_j \phi^+) + H.c. \tag{1}$$

induces Majorana neutrino masses as the Higgs scalar ϕ^0 acquires a nonzero vacuum expectation value $\langle \phi^0 \rangle = v$, so that

$$(\mathcal{M}_\nu)_{ij} = \frac{f_{ij} v^2}{\Lambda}. \tag{2}$$

This shows that neutrino mass is seesaw in character, i.e. it is inversely proportional to some large scale Λ. The ultraviolet completion of this effective operator may be accomplished in three ways at tree level [2] using (I) heavy Majorana fermion singlets N_i, (II) a heavy scalar triplet (ξ^{++}, ξ^+, ξ^0), or (III) heavy Majorana fermion triplets $(\Sigma^+, \Sigma^0, \Sigma^0)_i$, commonly referred to as Type I, Type II, or Type III seesaw. There are also three one-particle-irreducible (1PI) one-loop realizations [2]. Recently the one-particle-reducible (1PR) diagrams have also been considered [3].

If there are new particles with masses below the elctroweak scale, such as fermion singlets ν_S with mass m_S, then neutrinos may acquire mass through their mixing with ν_S. However, this mechanism is still seesaw because m_ν is still inversely proportional to m_S. There is however an exception. It has been pointed out recently [4] that in the scotogenic model of radiative neutrino mass [5], it is possible to have m_ν directly proportional to m_S, and there is no mixing between m_ν and m_S.

This model was proposed [5] in 2006 to connect neutrino mass with dark matter. The idea is very simple. Assume three neutral fermion singlets N_i as in the usual Type I seesaw [6], but let them be odd under a new Z_2 symmetry, so that there is no $(\nu_i \phi^0 - l_i \phi^+)N_j$ coupling and the effective operator of Eq. (1) is not realized. At this stage, N_i may have Majorana masses M_i, but ν_i is massless. However, they can be linked through the interaction $h_{ij}(\nu_i\eta^0 - l_i\eta^+)N_j$ where (η^+, η^0) is a new scalar doublet which is also odd under the aforementioned Z_2 [7]. Hence Majorana neutrino masses are generated in one loop as shown in Fig. 1.

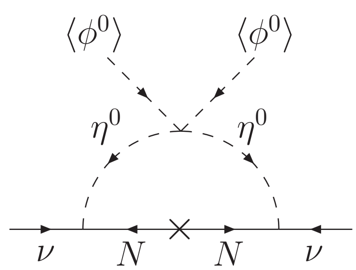

Figure 1: One-loop generation of scotogenic Majorana neutrino mass.

This mechanism has been called "scotogenic", from the Greek "scotos" meaning darkness. Because of the allowed $(\lambda_5/2)(\Phi^\dagger\eta)^2 + H.c.$ interaction, $\eta^0 = (\eta_R + i\eta_I)/\sqrt{2}$ is split so that $m_R \neq m_I$. The diagram of Fig. 1 can be computed exactly [5], i.e.

$$(\mathcal{M}_\nu)_{ij} = \sum_k \frac{h_{ik}h_{jk}M_k}{16\pi^2}\left[\frac{m_R^2}{m_R^2 - M_k^2}\ln\frac{m_R^2}{M_k^2} - \frac{m_I^2}{m_I^2 - M_k^2}\ln\frac{m_I^2}{M_k^2}\right]. \tag{3}$$

A good dark-matter candidate is η_R as first pointed out in Ref. [5]. It was subsequently proposed by itself in Ref. [8] and studied in detail in Ref. [9]. The η doublet has become known as the "inert" Higgs doublet, but it does have gauge and scalar interactions even if it is the sole addition to the standard model.

The usual assumption for neutrino mass in Eq. (1) is

$$m_I^2 - m_R^2 << m_I^2 + m_R^2 << M_k^2, \tag{4}$$

in which case

$$(\mathcal{M}_\nu)_{ij} = \frac{\lambda_5 v^2}{8\pi^2}\sum_k \frac{h_{ik}h_{jk}}{M_k}\left[\ln\frac{M_k^2}{m_0^2} - 1\right], \tag{5}$$

where $m_0^2 = (m_I^2 + m_R^2)/2$ and $m_R^2 - m_I^2 = 2\lambda_5 v^2$ $(v = \langle \phi^0 \rangle)$. This scenario is often referred to as the radiative seesaw. What was not realized in most applications of this model since 2006 is that there is another very interesting scenario, i.e.

$$M_k^2 << m_R^2,\ m_I^2. \tag{6}$$

Neutrino masses are then given by [4]

$$(\mathcal{M}_\nu)_{ij} = \frac{\ln(m_R^2/m_I^2)}{16\pi^2} \sum_k h_{ik} h_{jk} M_k. \tag{7}$$

This simple expression is actually very extraordinary, because neutrino mass is now not inversely proportional to some large scale. In that case, how do we understand the smallness of m_ν? The answer is lepton number. In this model, $(\nu, l)_i$ have lepton number $L = 1$ and N_k have $L = -1$, and L is conserved in all interactions except for the Majorana mass terms M_k which break L to $(-1)^L$. We may thus argue that M_k should be small compared to all other mass terms which conserve L, the smallest of which is the electron mass, $m_e = 0.511$ MeV. It is thus reasonable to have $M_k \sim 10$ keV, in which case $m_\nu \sim 0.1$ eV is obtained if $h^2 \sim 10^{-3}$ in Eq. (7). Each neutrino mass is then simply proportional to a linear combination of M_k according to Eq. (7). Their ratio is just a scale factor and small neutrino masses are due to this "scaling" mechanism. Note that the interesting special case where only M_1 is small has been considered previously [10, 11]. Note also that if $|m_I^2 - m_R^2| = 2|\lambda_5|v^2 << |m_I^2 + m_R^2|$, then $\ln(m_R^2/m_I^2)$ would be strongly suppressed, but this is not compulsory. For example, let $m_R = 240$ GeV, $m_I = 150$ GeV, then $|\lambda_5| = 0.58$ and $\ln(m_R^2/m_I^2) = 0.94$.

The scotogenic model [5] with large M_k, i.e. Eq. (5), has been extended recently [12] to include the well-known non-Abelian discrete symmetry A_4 [13, 14, 15]. Here we consider the case of Eq. (7). This assumption changes the phenomenology of N_k as well as (η^+, η^0) and may render this model to be more easily verifiable at the Large Hadron Collider (LHC). We let (η^+, η^0) be a singlet under A_4 and both (ν_i, l_i) and N_k to be triplets. In that case,

$$h_{ik} = h\delta_{ik}, \tag{8}$$

and

$$\mathcal{M}_\nu = \zeta \mathcal{M}_N, \tag{9}$$

where $\zeta = h^2 \ln(m_R^2/m_I^2)/16\pi^2$ is the scale factor. The soft breaking of A_4 which shapes $\mathcal{M}_N$ is then directly transmitted to $\mathcal{M}_\nu$.

One immediate consequence of this restricted scaling mechanism for neutrino mass is that if $M_{1,2,3}$ are all of order 10 keV, then the three neutrino masses are all of order 0.1 eV, i.e. a quasidegenerate scenario. For example, if $m_1 = 0.1$ eV and is the lightest, then for $M_1 = 10$ keV, $M_3 = 10^5(m_1 + \sqrt{\Delta m^2_{31}}) = 14.85$ keV. Another immediate consequence is that the interactions of $N_{1,2,3}$ with the charged leptons through η^+ depend only on h and the mismatch between the charged-lepton mass matrix and the neutrino mass matrix, i.e. the experimentally determined neutrino mixing matrix $U_{l\nu}$. Hence $\mu \to e\gamma$ is highly suppressed because the leading term of its amplitude is proportional to $\sum_k h_{\mu k} h^*_{ek} = |h|^2 \sum_k U_{\mu k} U^*_{ek} = 0$. The next term $\sum_k U_{\mu k} U^*_{ek} M_k^2/m^2_{\eta^+}$ is nonzero but is negligibly small. This means that there is no useful bound on the η^+ mass from $\mu \to e\gamma$. *Note that A_4 may be replaced by any other flavor symmetry as long as it is possible to have Eq. (8) using the singlet and triplet representations of that symmetry.* As for the muon anomalous magnetic moment, it is given by [16]

$$\Delta a_\mu = -\frac{m_\mu^2 |h|^2}{96\pi^2 m^2_{\eta^+}} = -1.18 \times 10^{-12} \left(\frac{|h|^2}{10^{-3}}\right) \left(\frac{100 \text{ GeV}}{m_{\eta^+}}\right)^2. \tag{10}$$

Since the experimental uncertaintly is 6×10^{-10}, this also does not give any useful bound on the η^+ mass.

If $\eta^\pm, \eta_R, \eta_I$ are of order 10^2 GeV, the interactions of N_k with the neutrinos and charged leptons are weaker than the usual weak interaction, hence N_k may be considered "sterile" and become excellent warm dark-matter candidates [17, 18]. However, unlike the usual sterile neutrinos [19] which mix with the active neutrinos, the lightest N_k here is absolutely stable. This removes one of the most stringent astrophysical constraints on warm dark matter, i.e. the absence of galactic X-ray emission from its decay, which would put an upper bound of perhaps 2.2 keV on its mass [20], whereas Lyman-α forest observations (which still apply in this case) impose a lower bound of perhaps 5.6 keV [21]. Such a stable sterile neutrino (called a "scotino") is also possible in an unusual left-right extension [22] of the standard model. Conventional left-right models where the $SU(2)_R$ neutrinos mix with the $SU(2)_L$ neutrinos have also been studied [23, 24, 25].

Since N_k are assumed light, muon decay proceeds at tree level through η^+ exchange, i.e. $\mu \to N_\mu e \bar{N}_e$. The inclusive rate is easily calculated to be

$$\Gamma(\mu \to N_\mu e \bar{N}_e) = \frac{|h|^4 m_\mu^5}{6144 \pi^3 m_{\eta^+}^4}. \tag{11}$$

Since N_μ and $\bar{N}_e$ are invisible just as ν_μ and $\bar{\nu}_e$ are invisible in the dominant decay $\mu \to \nu_\mu e \bar{\nu}_e$ (with rate $G_F^2 m_\mu^5 / 192\pi^3$), this would change the experimental value of G_F. Using the experimental uncertainty of 10^{-5} in the determination of G_F, we find

$$m_{\eta^+} > 70 \text{ GeV} \tag{12}$$

for $|h|^2 = 10^{-3}$. This is a useful bound on the η^+ mass, but it is also small enough so that η^+ may be observable at the LHC. The phenomenological bound on m_{η^+} from e^+e^- production at LEPII has been estimated [26] to be $70 - 90$ GeV. A bound of 80 GeV was used in a previous study [27] of this model.

Whereas the lightest scotino, say N_1, is absolutely stable, $N_{2,3}$ will decay into N_1 through η_R and η_I. The decay rate of $N_3 \to N_1 \bar{\nu}_1 \nu_3$ is given by

$$\Gamma(N_3 \to N_1 \bar{\nu}_1 \nu_3) = \frac{|h|^4}{256 \pi^3 M_3} \left(\frac{1}{m_R^2} + \frac{1}{m_I^2} \right)^2$$

$$\times \left(\frac{M_3^6}{96} - \frac{M_1^2 M_3^4}{12} + \frac{M_1^6}{12} - \frac{M_1^8}{96 M_3^2} + \frac{M_1^4 M_3^2}{8} \ln \frac{M_3^2}{M_1^2} \right). \tag{13}$$

Let $M_1 = 10$ keV, $M_3 = 14.85$ keV, $|h|^2 = 10^{-3}$, $m_R = 240$ GeV, $m_I = 150$ GeV, then this rate is 1.0×10^{-46} GeV, corresponding to a lifetime of 2.1×10^{14} y, which is much longer than the age of the Universe of $13.75 \pm 0.11 \times 10^9$ y. The lifetime of N_2 is even longer because $\Delta m^2_{21} << \Delta m^2_{31}$. Hence both N_2 and N_3 are stable enough to be components of warm dark matter. However, $N_{2,3} \to N_1\gamma$ are negligible for the same reason that $\mu \to e\gamma$ is negligible, so they again have no galactic X-ray signatures.

$m_{\eta^\pm}$ (GeV)	$\not{E}_T > 0$ GeV	$\not{E}_T > 25$ GeV	$\not{E}_T > 50$ GeV	$\not{E}_T > 100$ GeV
80	33.2	27.9	18.3	2.88
90	22.7	19.8	14.4	3.10
100	15.7	14.0	10.6	3.08
110	11.4	10.3	8.13	3.03
120	8.72	7.99	6.54	2.91
130	6.45	5.98	5.05	2.57
140	4.97	4.64	3.96	2.21
150	3.84	3.62	3.16	1.89
SM Background	626.45	453.3	205.8	8.6

Table 1: Cross sections (fb) for signal and dominant SM background at LHC with E_{CM}= 8 TeV. For the signal, $pp \to \eta^\pm\eta^\mp \to e^\pm\mu^\mp N_1 N_2$ (fb) with various $\not{E}_T$ cuts and different m_η (GeV) are specified. For SM background, W^+W^- production and decays to $e^\pm\mu^\mp$+ $\not{E}_T$ with the same cuts are specified.

Since η^+ may be as light as 70 GeV, it may be observable at the LHC. Assuming that the recently observed particle [28, 29] at the LHC is the Higgs boson H coming from (ϕ^+, ϕ^0), the decay $H \to \eta^+\eta^-$ is not allowed for $m_H = 126$ GeV. However, $\eta^\pm$ will contribute to the $H \to \gamma\gamma$ rate, as already pointed out [30, 31, 32, 33]. What sets our model apart is the inclusive decay of $\eta^\pm \to l^\pm N_{1,2,3}$, which is of universal strength. At the LHC, the pair production of $\eta^+\eta^-$ will then lead to $l_i^+ l_j^-$ final states with equal probability for each flavor combination. For example, $e^+\mu^-$ and μ^+e^- will each occur 1/9 of the time. This signature together with the large missing energy of $N_{1,2,3}$ may allow it to be observed at the LHC. However, these events also come from W^+W^- production and their subsequent leptonic decays. As shown in Table 1, the cross sections for the signal events are smaller than the dominant W^+W^- background even after a large $\not{E}_T$ cut. This is because both the signal and background events have similar missing energy distributions, but the W^+W^- production is much larger. If data show an excess of such events [34] over the SM prediction, it could be due to $\eta^+\eta^-$, but it may also simply come from an incorrect scale factor used in the SM calculation. Hence it is difficult to draw any conclusion for the case at hand. We generated the signal events using `Calchep` [35] and interfacing them to the event generator `Pythia` [36]. Background events were generated with `Pythia` and an appropriate $K-$factor was applied to match the NLO cross-section for W^+W^- at 8 TeV for the LHC, which is 57.3 pb [34]. We used CTEQ6L [37] parton distribution functions for both signal and background events. The basic $p_T > 10$ GeV and $|\eta| < 2.5$ cuts are applied for leptons.

In the supersymmetric $SU(5)$ completion [38] of this model, there are exotic quarks which may be produced abundantly. Their decays into $\eta^\pm$ would have four leptons of different flavor in the final state. This may be a better signature of this model. Details will be given elsewhere.

In conclusion, the scotogenic model [5] of neutrino mass with a solution [4] where there is no seesaw mechanism and $N_{1,2,3}$ have masses of order 10 keV has been implemented with the non-Abelian discrete symmetry A_4. The scotinos $N_{1,2,3}$ are good warm dark-matter candidates which can explain the structure of the Universe at all scales [17, 18]. Since N_1 is absolutely stable and the decays $N_{2,3} \to N_1\gamma$ are negligible, the galactic X-ray upper bound of perhaps 2.2 keV on its mass [20] is avoided. It will also not be detected in terrestrial experiments. On the other hand, since this model requires an extra scalar doublet, and $\eta^\pm$

may be as light as 70 GeV, it may be tested at the LHC, especially if it is the decay product of an exotic quark.

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