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Radiative Scaling Neutrino Mass with A_4 Symmetry and Warm Dark Matter

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Abstract

In scotogenic model, neutrino mass comes from a one-loop mechanism with dark matter candidate in the loop. An assumption of producing small neutrino mass proportional to small dark matter mass in the model consisting of singlet Majorana fermions N_i with masses of order 10 keV was proposed recently. Neutrino masses are scaled down from them by factors of about 10^{-5}. We study the consequences in the phenomenology of N_i as warm dark matter of implementing the non-Abelian discrete symmetry A_4 for neutrino mixing in the model.

In the scotogenic model [1] the radiative neutrino mass connects with dark matter with considering three neutral fermion singlets N_i and a new scalar doublet (η^+, η^0) are odd under a new Z_2 symmetry. Hence Majorana neutrino masses are generated in one loop through the interaction $h_{ij}(\nu_i \eta^0 - l_i \eta^+) N_j$ as shown in Fig. 1. 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 [1], i.e.

$$(\mathcal{M}_\nu)_{ij} = \sum_k \frac{h_{ik} h_{jk} M_k}{16\pi^2} \left[\frac{m_R^2}{m_R^2 - M_k^2} \ln \frac{m_R^2}{M_k^2} - \frac{m_I^2}{m_I^2 - M_k^2} \ln \frac{m_I^2}{M_k^2} \right]. \quad (1)$$

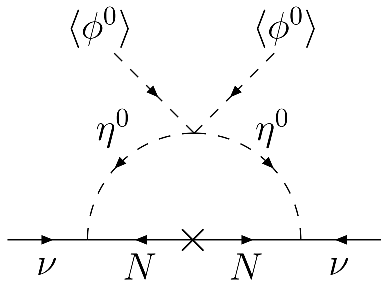

Figure 1: One-loop generation of scotogenic Majorana neutrino mass.

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

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

in which case

$$(\mathcal{M}_\nu)_{ij} = \frac{\lambda_5 v^2}{8\pi^2} \sum_k \frac{h_{ik} h_{jk}}{M_k} \left[\ln \frac{M_k^2}{m_0^2} - 1 \right], \tag{3}$$

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. There is another interesting scenario, i.e.

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

Neutrino masses are then given by [2]

$$(\mathcal{M}_\nu)_{ij} = \frac{\ln(m_R^2/m_I^2)}{16\pi^2} \sum_k h_{ik} h_{jk} M_k. \tag{5}$$

In the above expression 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. For instance, if $M_k \sim 10$ keV then $m_\nu \sim 0.1$ eV is obtained if $h^2 \sim 10^{-3}$ in Eq. (5).

The scotogenic model [1] with large M_k, i.e. Eq. (3), has been extended recently [3] to include the well-known non-Abelian discrete symmetry A_4 [4, 5, 6]. Simpler discrete symmetries, such as Z_2 [7, 8], have been used in the literature to describe the neutrino mixing. In Eq. (5) we let (η^+, η^0) be a singlet under A_4 and both (ν_i, l_i) and N_k to be triplets [9]. In that case ,

$$h_{ik} = h\delta_{ik}, \tag{6}$$

and

$$\mathcal{M}_\nu = \zeta \mathcal{M}_N, \tag{7}$$

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 $\mathcal{M}_\nu$.

There are two consequences of the scaling mechanism for neutrino mass. First, quasidegenerate scenario of neutrino mass can be obtained if $M_{1,2,3}$ are in the same order of magnitude. Second, 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 $\sum_k h_{\mu k} h^*_{ek} = |h|^2 \sum_k U_{\mu k} U^*_{ek} = 0$. *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.*

N_k becomes good warm dark-matter candidates [10, 11] if the interactions of N_k with the neutrinos and charged leptons are weaker than the usual weak interaction when $\eta^\pm, \eta_R, \eta_I$ are of order 10^2 GeV, i.e. N_k may be considered "sterile". Since N_k are assumed light in the range of keV, muon decay proceeds at tree level through η^+ exchange, i.e. $\mu \to N_\mu e \bar{N}_e$, with the rate

$$\Gamma(\mu \to N_\mu e \bar{N}_e) = \frac{|h|^4 m_\mu^5}{6144\pi^3 m_{\eta^+}^4}. \tag{8}$$

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

$$m_{\eta^+} > 70 \text{ GeV} \tag{9}$$

for $|h|^2 = 10^{-3}$. 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

$$\begin{aligned}\Gamma(N_3 \to N_1 \bar{\nu}_1 \nu_3) &= \frac{|h|^4}{256\pi^3 M_3} \left(\frac{1}{m_R^2} + \frac{1}{m_I^2} \right)^2 \\ &\times \left(\frac{M_3^6}{96} - \frac{M_1^2 M_3^4}{12} + \frac{M_1^6}{12} - \frac{M_1^8}{96 M_3^2} + \frac{M_1^4 M_3^2}{8} \ln \frac{M_3^2}{M_1^2} \right). \end{aligned} \tag{10}$$

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 $\Delta m_{21}^2 << \Delta m_{31}^2$. Hence both N_2 and N_3 are stable enough to be components of warm dark matter.

Since η^+ may be as light as 70 GeV, it may be observable at the LHC. 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. 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. 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.

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