

Neutrino dark energy and leptogenesis with TeV scale triplets

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Received: 2 March 2016 / Accepted: 14 April 2016 / Published online: 28 April 2016

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Abstract We propose a realization of mass varying neutrino dark energy in two extensions of the Standard Model (SM) with a dynamical neutrino mass related to the acceleron field while satisfying the naturalness. In the first scenario the SM is extended to include a TeV scale scalar Higgs triplet (ξ) and a TeV scale second Higgs doublet (η), while in the second scenario an extension of the SM with fermion triplet (Σ) is considered. We also point out the possible leptogenesis mechanisms for simultaneously generating the observed baryon asymmetry of the universe in both scenarios and discuss the collider signatures for the TeV scale new fields which make these models testable in the current run of LHC.

1 Introduction

Evidence from the astrophysical observations suggests that out of the total mass–energy budget of the universe, the baryonic and dark matter together account for only about 30 % while the remaining 70 %, referred to as dark energy, is attributed to the accelerated expansion of the universe and remains a challenge to explain. While the existence of a scalar field called quintessence provides an explanation, the striking proximity of the effective scales of neutrino masses and the dark energy points to a connection between them, realized in the neutrino dark energy (ν DE) models. To this end, several approaches have been used in the literature [1–24]. In some scenarios, a direct connection through the formation of a neutrino condensate at a late epoch of the early universe using the effective self-interaction has been explored [1–7]. Another class of models invoke the variation of neutrino masses to dynamically obtain the dark energy [8–23]. In this work, we will focus on the latter approach.

The atmospheric, solar, and reactor neutrino oscillation experiments have confirmed the existence of tiny but non-

zero masses of neutrinos ($\sim 10^{-2}$ eV). An attractive explanation of the neutrino masses employs the seesaw mechanism [25–29], giving rise to naturally small Majorana or Dirac masses of neutrinos. In addition, the baryon asymmetry of the universe can be generated through leptogenesis [30] in the framework of the seesaw scenario. In the original ν DE models, the Standard Model (SM) is extended to accommodate singlet right-handed neutrinos (N_i , $i = 1, 2, 3$) giving a Majorana mass to the light neutrinos. The Majorana masses of the right-handed neutrinos are made to vary with the acceleron field, connecting the light neutrino masses with the scale of dark energy. However, the naturalness requires the Majorana masses of the right-handed neutrinos to be in the eV range, in contradiction to the expected scenario of a very heavy M_{N_i} triggering the canonical seesaw mechanism. In Ref. [18], it was pointed out that the above problem can be avoided if the SM is extended to include triplet Higgs scalars. However, in such a scheme the coefficient of the trilinear scalar coupling with mass dimension varies with the acceleron field and this predicts the mass scale of the triplet Higgs scalars to be close to the electroweak symmetry breaking scale (of order 100 GeV), which has not been observed at the LHC so far.

The purpose of this work is to point out two ways to get around the above constraint, while simultaneously explaining the observed baryon asymmetry of the universe. One way is to add some additional scalar field to push the additional scalar field masses to TeV scale, readily testable at the current run of LHC. Another way is to add fermion triplets instead of scalar triplets and utilize the type III seesaw scheme.

First we propose a realization of mass varying neutrinos in an extension of the usual triplet Higgs model which includes a second Higgs doublet (η) in addition to the SM Higgs doublet (Φ) and Higgs triplet (ξ), but no right-handed neutrinos [31,32]. In this scenario both additional Higgs fields are of TeV scale and the smallness of the neutrino mass comes from the lepton number breaking scalar sector. This model has highly predictive collider signatures and thus it can be

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right away put to test in the current run of LHC. Next we propose a new model of ν DE utilizing an extension of the SM with fermion triplet $(\Sigma^+, \Sigma^0, \Sigma^-)_R$, where the neutrino mass is dynamical and related to the accelaron field. This model can naturally give the correct energy scale associated with the neutrino mass and it provides a rich TeV scale phenomenology, testable at the LHC. We also point out possible leptogenesis mechanisms for simultaneously generating the observed baryon asymmetry of the universe in both models.

2 Neutrino masses and dark energy connection

By extending the SM to include a heavy Higgs triplet (ξ^{++}, ξ^+, ξ^0) with trilinear couplings to both the lepton doublet $L_i = (v_i, l_i)$ and the Higgs doublet $\Phi = (\phi^+, \phi^0)$, one can realize the unique dimension-five effective operator [33]

$$\mathcal{L}_{\text{eff}} = \frac{f_{ij}}{\Lambda} L_i L_j \Phi \Phi, \quad (1)$$

obtained by integrating out the heavy degrees of freedom (with mass much larger than the ordinary SM particles) associated to a characteristic heavy mass scale Λ . Thus the neutrinos, massless in the minimal SM, acquire small Majorana masses. The relevant interaction terms are given by

$$\mathcal{L}_{\text{int}} = f_{ij} \left[v_i v_j \xi^0 + \frac{1}{\sqrt{2}} (v_i l_i + l_j v_j) \xi^+ + l_i l_j \xi^{++} \right] + \text{h.c.}, \quad (2)$$

where $\tilde{\Phi} = (\tilde{\phi}^0, -\phi^-)$. The above interaction terms give [31]

$$(M_\nu)_{ij} = \frac{2 f_{ij} \mu \langle \phi^0 \rangle^2}{m_{\xi^0}^2}. \quad (3)$$

Thus it follows that if μ is a function of the accelaron field \mathcal{A} i.e. $\mu = \mu(\mathcal{A})$, then the mass varying neutrinos can be realized for m_ξ of the order of the electroweak scale. However, if the ν DE is indeed realized through the Higgs triplet, then at least ξ^{++} should have been observable at the LHC. Thus it is worth exploring if such a Higgs triplet can be schemed to have a mass of TeV scale in light of the current run of LHC.

2.1 Model A

In the presence of the additional Higgs doublet η in the above scheme, the neutrino masses come from the Higgs triplet ξ (with lepton number assignment $L = -2$) and its interaction with η (carrying lepton number $L = -1$) [32]. The most general lepton number conserving scalar potential is given by

$$\begin{aligned} V = & m_1^2 \Phi^\dagger \Phi + m_2^2 \eta^\dagger \eta + m_3^2 \text{Tr}[\Delta^\dagger \Delta] \\ & + \frac{1}{2} \Lambda_1 (\Phi^\dagger \Phi)^2 + \frac{1}{2} \Lambda_2 (\eta^\dagger \eta)^2 + \frac{1}{2} \Lambda_3 (\text{Tr}[\Delta^\dagger \Delta])^2 \\ & + \frac{1}{2} \Lambda_4 (\text{Tr}[\Delta^\dagger \Delta^\dagger]) (\text{Tr}[\Delta \Delta]) + \Lambda_5 (\Phi^\dagger \Phi) (\eta^\dagger \eta) \\ & + \Lambda_6 (\Phi^\dagger \Phi) (\text{Tr}[\Delta^\dagger \Delta]) + \Lambda_7 (\eta^\dagger \eta) (\text{Tr}[\Delta^\dagger \Delta]) \\ & + \Lambda_8 (\Phi^\dagger \eta) (\eta^\dagger \Phi) + \Lambda_9 (\Phi^\dagger \Delta^\dagger \Delta \Phi) + \Lambda_{10} (\eta^\dagger \Delta^\dagger \Delta \eta) \\ & + s \mu (\eta^\dagger \Delta \tilde{\eta}) + \text{h.c.}, \end{aligned} \quad (4)$$

where

$$\Delta = \begin{pmatrix} \xi^+/\sqrt{2} & \xi^{++} \\ \xi^0 & -\xi^+/\sqrt{2} \end{pmatrix}, \quad (5)$$

$\tilde{\eta} = (\tilde{\eta}^0, -\eta^-)$ and μ has the dimension of mass. The lepton number is softly broken by the terms

$$V_{\text{soft}} = \mu_1^2 \Phi^\dagger \eta + \mu_2 (\Phi^\dagger \Delta \tilde{\eta}) + \mu_3 (\Phi^\dagger \Delta \tilde{\Phi}) + \text{h.c.} \quad (6)$$

Next we define vacuum expectation values (VEVs) of the scalar fields to be $\langle \phi^0 \rangle = v_1$, $\langle \eta^0 \rangle = v_2$, and $\langle \xi^0 \rangle = v_3$. Now minimization of the potential with respect to the various Higgs fields give the consistency conditions and the relations between the different VEVs, which can be solved assuming $m_1^2 < 0$, but $m_2^2 > 0$ and $m_3^2 > 0$ to obtain

$$\begin{aligned} v_1^2 & \simeq -m_1^2/\Lambda_1, \\ v_2 & \simeq -\mu_1^2 v_1 / [m_2^2 + (\Lambda_5 + \Lambda_8) v_1^2], \\ v_3 & \simeq -(\mu v_2^2 + \mu_2 v_1 v_2 + \mu_3 v_1^2) / (m_3^2 + \Lambda_6 v_1^2). \end{aligned} \quad (7)$$

Thus taking m_2, m_3 , and μ to be $M \sim \text{TeV}$ we have

$$v_2 \sim \mu_1^2 v_1 / M^2, \quad v_3 \sim v_2^2 / M. \quad (8)$$

Consequently, $u \ll v_2 \ll v_1$ and

$$v_3 \sim \mu_1^2 v_1^2 / M^5. \quad (9)$$

For $v_1 \sim 10^2 \text{ GeV}$ and $\mu_1 \sim 1 \text{ GeV}$ we have $v_2 \sim 0.1 \text{ MeV}$ and $v_3 \sim 10^{-2} \text{ eV}$, which gives the correct order of magnitude for neutrino mass $(m_\nu)_{ij} = 2 f_{ij} v_3$. Thus we have a natural realization of the required small neutrino masses with TeV scale additional Higgs fields, which does not need any large extra space dimensions constraining m_ξ below the cutoff energy scale. Moreover, this model is much more flexible compared to the scenario with only Higgs triplet in the sense that there is no strict constraint on m_ξ to be of the order of electroweak scale. Now the realization of the ν DE model through mass varying neutrinos is straightforward. The idea is to make μ_1 a function of the accelaron field \mathcal{A} , i.e. $\mu_1 = \mu_1(\mathcal{A})$. We will come back to the realization of ν DE once we give the account of the other model below.

2.2 Model B

The extension of the fermion (lepton) sector of the SM can be realized in two ways. The basic idea is that the new lepton multiplet gains a large mass and then it mixes with the ordinary lepton doublet triggering the seesaw mechanism. The new lepton multiplet can only be a singlet or a triplet of $SU(2)_L$. The idea of the triplet lepton representation to utilize the seesaw structure in neutrino mass matrix was first proposed in Ref. [34], referred to as type III seesaw, which have been generalized in the context of unified theories in Ref. [35]. The simplest way to utilize the type III seesaw is to add the $SU(2)_L$ triplet with zero hypercharge,

$$\Sigma = \begin{pmatrix} \Sigma^0/\sqrt{2} & \Sigma^+ \\ \Sigma^- & -\Sigma^0/\sqrt{2} \end{pmatrix}, \quad (10)$$

to the SM, with the interaction terms $\bar{\Sigma}$

$$\mathcal{L}_{\text{int}} = -\frac{1}{2} \text{Tr} [\bar{\Sigma} M_{\Sigma} \Sigma^c + \bar{\Sigma}^c M_{\Sigma}^* \Sigma] - \tilde{\Phi}^\dagger \bar{\Sigma} \sqrt{2} Y_{\Sigma} L - \bar{L} \sqrt{2} Y_{\Sigma}^\dagger \Sigma \tilde{\Phi}. \quad (11)$$

The terms related to neutrino mass matrix can be identified readily to obtain the mass matrix:

$$\mathcal{L}_{\nu, \text{mass}} = (\nu \ \Sigma^0) \begin{pmatrix} 0 & Y_{\Sigma} v/2\sqrt{2} \\ Y_{\Sigma}^T v/2\sqrt{2} & M_{\Sigma}/2 \end{pmatrix} \begin{pmatrix} \nu \\ \Sigma^0 \end{pmatrix}. \quad (12)$$

This gives the non-zero neutrino masses given by

$$M_{\nu} = -\frac{v^2}{2} Y_{\Sigma}^T M_{\Sigma}^{-1} Y_{\Sigma}. \quad (13)$$

Now the neutrino masses can be connected to the dark energy by simply taking $M_{\Sigma} = M_{\Sigma}(\mathcal{A})$; however, such a scenario is constrained from naturalness. This scheme can be generalized right away by accommodating the right-handed neutrinos N_i^c , $i = 1, 2, 3$ in the scenario. The most general neutrino mass matrix in such a scenario can be written as

$$\mathcal{L}_{\nu, \text{mass}} = (\nu \ N^c \ \Sigma^0) \begin{pmatrix} 0 & M_N & F_1 u \\ M_N^T & 0 & F_2 \Omega \\ F_1^T u & F_2^T \Omega & M_{\Sigma} \end{pmatrix} \begin{pmatrix} \nu \\ N^c \\ \Sigma^0 \end{pmatrix}, \quad (14)$$

where the off-diagonal terms in the third column and row correspond to the mass terms $F_{1ij} \nu_i \Sigma^0 u$ and $F_{2ij} N_i^c \Sigma^0 \Omega$ with u and Ω being the VEVs of the corresponding Higgs fields. The realization of the nonrenormalizable term, giving rise to the type III seesaw, by integrating out heavy fields, is shown in Fig. 1.

The neutrino mass in the above scenario has two kinds of contributions, given by

$$M_{\nu} = -M_N (F_2 \Omega M_{\Sigma}^{-1} F_2^T \Omega)^{-1} M_N^T - (M_N + M_N^T) \frac{u}{\Omega}, \quad (15)$$

where the first term corresponds to a “double seesaw” contribution and the second term corresponds to the type III seesaw

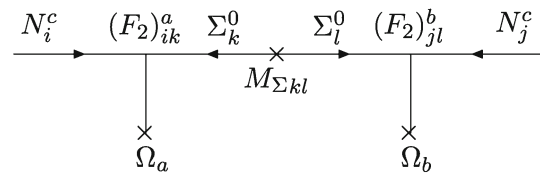


Fig. 1 Diagram realizing the effective nonrenormalizable operator generating the right-handed neutrino mass $M_R = (F_2 \Omega) M_{\Sigma}^{-1} (F_2^T \Omega)$

contribution. The relative contributions of the two kinds of terms to M_{ν} is model dependent. To satisfy the naturalness constraints we will consider the case $M_N \sim 1$ eV. Now taking $u \sim v \sim 10^2$ GeV, $\Omega \sim 10^4$ GeV and considering the phenomenologically interesting case $M_{\Sigma} \sim 10^3$ GeV with verifiable implications at the current run of LHC, it follows that for $F_2 \gtrsim 10^{-6}$ the dominant contribution to M_{ν} in Eq. (15) comes from the second term associated with the type III seesaw contribution and for the above set of values we obtain $M_{\nu} \sim 10^{-2}$ eV as desired. The mass varying neutrinos can be realized by taking $M_N = M_N(\mathcal{A})$.

2.3 Realization of neutrino dark energy

Having given the details of the two models realizing mass varying neutrinos with desired small masses while satisfying the naturalness, we are now ready to discuss the realization of ν DE where the neutrino mass (assumed to be a function of the canonically normalized acceleration field \mathcal{A}) $M_{\nu}(\mathcal{A})$ is a dynamical quantity and $\partial M_{\nu}/\partial \mathcal{A} \neq 0$ [9]. In the nonrelativistic limit, the energy density consists of the thermal neutrino (and antineutrino) background ($M_{\nu} n_{\nu}$) and the scalar potential $V_0(M_{\nu})$. The effective potential can be written as

$$V(M_{\nu}) = M_{\nu} n_{\nu} + V_0(M_{\nu}). \quad (16)$$

The neutrino background (driving M_{ν} to small values) gets diluted as the universe expands and the source term decreases as a result, while V_0 is minimized for a large M_{ν} . Thus the two terms act in the opposite directions with a minimum at some intermediate M_{ν} with a non-zero V_0 . The minimum of the effective potential is given by

$$V'(M_{\nu}) = n_{\nu} + V'_0(M_{\nu}) = 0. \quad (17)$$

Now at any instant of time assuming the simple equation of state

$$p(t) = \Omega \rho(t), \quad (18)$$

it follows that

$$\Omega + 1 = -\frac{\partial \log V}{3 \partial \log a} = -\frac{M_{\nu} V'_0(M_{\nu})}{V} = \frac{\Omega_{\nu}}{\Omega_{\nu} + \Omega_{\mathcal{A}}}, \quad (19)$$

where $\Omega_{\nu} = M_{\nu} n_{\nu}/\rho_c$ is the neutrino energy density and $\Omega_{\mathcal{A}} = \rho_{\mathcal{A}}/\rho_c$ corresponds to the contribution of $V_0(M_{\nu})$ to the energy density, with ρ_c is the critical density and a is the

cosmic scale factor. Since the observed value of $\Omega \simeq -1$, Eq. (19) implies that the energy density in the thermal neutrino background must be much less compared to the total dark energy density. This in turn suggests that the potential $V_0(M_\nu)$ should be a flat potential. For the case where $d\Omega/dn_\nu$ is small, the relation

$$M_\nu \propto n_\nu^\Omega \quad (20)$$

holds. The above considerations are independent from any specific model of the neutrino mass [9] and we will use them to draw out the phenomenological consequences specific to the two models of interest.

As we have discussed above in model A, $\mu_1 = \mu_1(\mathcal{A})$ makes the effective mass of the neutrinos to vary. While in model B, $M_N = M_N(\mathcal{A})$ does the same. Now for the self interactions of the acceleron field \mathcal{A} we take the effective potential of the form

$$V_0 = \Lambda^4 \log(1 + |\mathcal{M}(\mathcal{A})/\bar{\mathcal{M}}|), \quad (21)$$

where in model A, $\mathcal{M}(\mathcal{A}) = \mu_1(\mathcal{A})$ and in model B, $\mathcal{M}(\mathcal{A}) = M_N(\mathcal{A})$. Hence Eq. (16) takes the form

$$V(x) = a_1 x + a_2 \log\left(1 + \frac{x}{a_3}\right), \quad (22)$$

where $x = M_\nu \propto |\mu(\mathcal{A})|$ and a_1, a_2, a_3 are all positive. Now assuming $\mathcal{M}(\mathcal{A})/\bar{\mathcal{M}} \gg 1$ it follows that $x_{\min} \propto a_2/a_1$ implying

$$M_\nu \propto n_\nu^{-1}, \quad (23)$$

which gives the desired $\Omega \simeq -1$. Thus, the two models under consideration can naturally explain the ν DE for the TeV scale ξ , the η masses in model A, and the TeV scale mass of the new fermion triplet Σ in model B. The TeV scale mass of these particles makes these models particularly interesting in the context of collider phenomenology at the LHC. We will come back to the implications and signatures of these two models for colliders such as the LHC, once we address the issue of leptogenesis in these two models.

3 Leptogenesis

In model A, the SM is extended to include scalar triplet and an additional Higgs doublet η , providing an attractive possibility of realizing a successful leptogenesis scenario. We start with the conventional formalism of scalar triplet leptogenesis in a hierarchical case. $SU(2)_L \times U(1)_Y$ is the valid gauge group at an energy scale far above the electroweak symmetry breaking. Thus it follows that if we analyze one of the three components of the triplet scalar field then the results will hold for the other two. From Eqs. (2), (4), and (6) we can read off the decay modes of ξ^{++} as

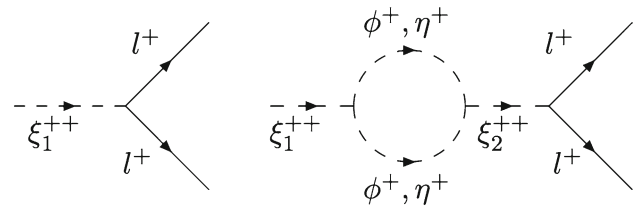


Fig. 2 The tree-level (left) and one-loop (right) decay diagrams for $\xi^{++} \rightarrow l^+ l^+$. A lepton asymmetry is generated by the CP violation occurring due to the interference between them

$$\xi_a^{++} \rightarrow \begin{cases} l_i^+ l_j^+ & (L = -2), \\ \phi^+ \phi^+ & (L = -2), \\ \eta^+ \eta^+ & (L = -0). \end{cases} \quad (24)$$

The coexistence of the above decay modes implies non-conservation of lepton number, however, the lepton asymmetry generated by ξ^{++} gets compensated by the decays of ξ^{--} , unless CP is also violated and the decays take place out of equilibrium. We follow the mass matrix formalism [31, 36], where the tree-level mass matrix for the triplets are assumed to be real and diagonal. Hence CP is conserved at tree level; however, CP conservation occurs at one-loop level due to interference between the tree and one-loop diagrams shown in Fig. 2.

Note that at least two ξ 's are required for CP nonconservation to occur. Following the mass-matrix formalism of Ref. [36], the diagonal tree-level mass matrix of ξ_a in Eq. (4) is modified in the presence of interactions to

$$\frac{1}{2} \xi^\dagger (M_+^2)_{ab} \xi_b + \frac{1}{2} (\xi_a^*)^\dagger (M_-^2)_{ab} \xi_b^*, \quad (25)$$

where

$$M_\pm^2 = \begin{pmatrix} M_1^2 - i\Gamma_{11} M_1 & -i\Gamma_{12}^\pm \\ -i\Gamma_{21}^\pm M_1 & M_2^2 - i\Gamma_{22} M_2 \end{pmatrix}, \quad (26)$$

with $\Gamma_{ab}^+ = \Gamma_{ab}$ and $\Gamma_{ab}^- = \Gamma_{ab}^*$. From the absorptive part of the one-loop diagram for $\xi_a \rightarrow \xi_b$ we have

$$\Gamma_{ab} M_b = \frac{1}{8\pi} \left(\mu_a \mu_b^* + \mu_{3a} \mu_{3b}^* + M_a M_b \sum_{k,l} f_{akl}^* f_{bkl} \right). \quad (27)$$

Now for $\Gamma_a \equiv \Gamma_{aa} \ll M_a$, the eigenvalues of M_\pm^2 are given by

$$\Lambda_{1,2} = \frac{1}{2} (M_1^2 + M_2^2 \pm \sqrt{S}), \quad (28)$$

where $S = (M_1^2 - M_2^2)^2 - 4|\Gamma_{12} M_2|^2$ and $M_1 > M_2$. The physical states are given by

$$\psi_{1,2}^+ = a_{1,2}^+ \xi_1 + b_{1,2}^+ \xi_2, \quad \psi_{1,2}^- = a_{1,2}^- \xi_1^* + b_{1,2}^- \xi_2^*, \quad (29)$$

where $a_1^\pm = b_2^\pm = 1/\sqrt{1 + |C_i^\pm|^2}$, $b_1^\pm = C_i^\pm/\sqrt{1 + |C_i^\pm|^2}$, $a_2^\pm = C_2^\pm/\sqrt{1 + |C_i^\pm|^2}$ with

$$C_1^+ = -C_2^- = \frac{-2i\Gamma_{12}^* M_2}{M_1^2 - M_2^2 + \sqrt{S}},$$

$$C_1^- = -C_2^+ = \frac{-2i\Gamma_{12} M_2}{M_1^2 - M_2^2 + \sqrt{S}}. \quad (30)$$

The states $\psi_{1,2}^\pm$ evolve with time and decay into lepton pair and antilepton pair.¹ Assuming $(M_1^2 - M_2^2)^2 \gg 4|\Gamma_{12} M_2|^2$, the lepton asymmetries generated are given by [31]

$$\varepsilon_i = \frac{1}{8\pi^2(M_1^2 - M_2^2)^2} \sum_{k,l} \{ \text{Im} [\mu_1 \mu_2^* f_{1kl} f_{2kl}^*] + \text{Im} [(\mu_3)_1 (\mu_3)_2^* f_{1kl} f_{2kl}^*] \} \left[\frac{M_i}{\Gamma_i} \right]. \quad (31)$$

For the case $M_1 > M_2$, when the temperature of the universe cools down below M_1 , ψ_1 decays away to create a lepton asymmetry. However, this asymmetry is washed out by lepton number nonconserving interactions of ψ_2 and the subsequent decay of ψ_2 at a temperature below M_2 sustains. The lepton asymmetry then gets converted to baryon asymmetry in the presence of the anomalous $B + L$ violating processes before the electroweak phase transition. The approximate final baryon asymmetry generated is given by

$$\frac{n_B}{s} \sim \frac{\varepsilon_2}{3g^* K (\ln K)^{0.6}}, \quad (32)$$

where $K \equiv \Gamma_2(M_2/T = 1)/H(M_2/T = 1)$ is a parameter measuring the deviation from thermal equilibrium, with the Hubble rate defined by $H = 1.66g^{*1/2}(T^2/M_{\text{Pl}})$, where g^* corresponds to the number of relativistic degrees of freedom.

Other than the decays and the inverse decays of triplet scalars, one needs to incorporate the gauge scatterings $\psi\bar{\psi} \leftrightarrow F\bar{F}, \phi\bar{\phi}, G\bar{G}$ (F corresponds to SM fermions and G corresponds to gauge bosons) and $\Delta L = 2$ scattering processes $ll \leftrightarrow \phi\phi^*$ and $l\phi \leftrightarrow \bar{l}\phi^*$ into the Boltzmann equation analysis of the asymmetry. Including the above washout processes, it turns out that $M_\xi \gtrsim 10^{11} \text{ GeV}$ is required in order to generate the correct asymmetry [38]. However, for a quasi-degenerate spectrum of scalar triplets the resonance effect can enhance the CP-asymmetry by a large amount and a successful leptogenesis scenario can be attained for a much smaller value of triplet scalar mass. A detailed analysis of the resonant leptogenesis is beyond the scope of this work and an account of the same can be found in Refs. [39,40], where an absolute bound of $M_\xi \gtrsim 1.6 \text{ TeV}$ is obtained for a successful leptogenesis scenario.

In model B, the type III seesaw scheme is realized and the right-handed neutrino mass matrix enters in the formula for

light neutrino masses compared to the type I seesaw. As a consequence, the light neutrino masses, mixing and leptogenesis are not that tightly coupled as in the case of the type I seesaw, where the constraints on the right-handed neutrino mass M_R can clash with the constraints coming from the textures of light neutrino masses and mixings. The advantage of the type III seesaw mechanism given in Eq. (14) is that, instead of the three heavy Majorana neutrinos in type I seesaw, here we have six heavy Majorana neutrinos. This can give rise to three pseudo-Dirac pairs of neutrinos with one or more pairs having degenerate masses. The six heavy two component neutrinos have the form of the mass matrix given by [41]

$$(\tilde{N}_i^c \quad \tilde{\Sigma}_i^0) \begin{pmatrix} 0 & M_i \delta_{ij} \\ M_i \delta_{ij} & \tilde{M}_{\Sigma ij} \end{pmatrix} \begin{pmatrix} \tilde{N}_j^c \\ \tilde{\Sigma}_j^0 \end{pmatrix}. \quad (33)$$

Now the degenerate lightest pair of pseudo-Dirac neutrinos or equivalently, two Majorana neutrinos $N_\pm \simeq (\tilde{N}_1^c \pm \tilde{\Sigma}_1^0)/\sqrt{2}$ with masses $M_\pm \simeq \pm M_1 + \frac{1}{2} \tilde{M}_{\Sigma 11}$ can decay into light neutrino and Higgs doublet via the Yukawa term $Y_{i\pm}(N_\pm \nu_i)\Phi$, where

$$Y_{i\pm} \simeq \frac{(\tilde{Y}_{i1} \pm (\tilde{F}_2)_{i1})}{\sqrt{2}} \mp \frac{\tilde{M}_{11}}{4M_1} \frac{(\tilde{Y}_{i1} \mp (\tilde{F}_2)_{i1})}{\sqrt{2}}. \quad (34)$$

The asymmetry generated by the decays of N_\pm is given by

$$\varepsilon_1 = \frac{1}{4\pi} \frac{\text{Im} [\sum_j (Y_{j+} Y_{j-}^*)]^2}{\sum_j (|Y_{j+}|^2 + |Y_{j-}|^2)} I(M_-^2/M_+^2), \quad (35)$$

where $I(M_-^2/M_+^2)$ comes from the absorptive part of the decay amplitude, with $I(x) = \sqrt{x}[1 - (1+x)\ln(1+(1/x)) + 1/(1-x)]$. Using the new basis parametrization $N_i^c = U_{ij} \tilde{N}_j^c$ and $\Sigma_{0i} = V_{ij} \tilde{\Sigma}_{0j}$ with the matrix $(F_1)_{ij}$ diagonal, where

$$U = \begin{pmatrix} u_{11} & \Lambda u_{12} & \Lambda u_{13} \\ \Lambda u_{21} & u_{22} & u_{23} \\ \Lambda u_{31} & u_{32} & u_{33} \end{pmatrix}, \quad (36)$$

with $u_{ij} \sim 1$ and

$$\tilde{F}_2 u = \begin{pmatrix} \Lambda^2 f_{11} & \Lambda f_{12} & \Lambda f_{13} \\ \Lambda f_{21} & f_{22} & f_{23} \\ \Lambda f_{31} & f_{32} & f_{33} \end{pmatrix} v_u, \quad (37)$$

the asymmetry can be put in the form [41]

$$\varepsilon_1 = \frac{\Lambda^2 (|u_{31}|^2 - |f_{31}|^2) \text{Im}(u_{31}^* f_{31})}{4\pi (|u_{31}|^2 + |f_{31}|^2 + |f_{21}|^2)} I. \quad (38)$$

The lepton asymmetry of the universe is computed using

$$Y_L = \frac{n_B}{s} \sim \frac{\varepsilon_2 d}{3g^* K (\ln K)^{0.6}}, \quad (39)$$

where d is the washout parameter. In this case, for a hierarchical mass spectrum of triplets the lower bound on the

¹ Note that ξ_a and ξ_a^* are CP conjugate states, while ψ_i^\pm are not.

triplet mass for a successful leptogenesis scenario is given by $M_\Sigma \gtrsim 3 \times 10^{10}$ [42, 43] and to have TeV scale leptogenesis one must assume a quasi-degenerate spectrum of fermion triplets giving resonant enhancement as in the case of scalar triplets, giving a TeV scale bound on M_Σ [39, 44].

4 Collider signatures

The triplet fields ξ and Σ can be produced at the LHC if their masses are of the order of TeV; therefore, LHC gives an unique opportunity to verify the mechanism of neutrino mass generation if any of these heavy states or their signatures are observed. To this end, we give a very brief summary of the production and observability of the triplet fields in the two models discussed above. A quantitative exploration of the discovery potential of these new fields is beyond the scope of this work and here we mainly concentrate on a qualitative account of the likely scenarios. The members of the scalar triplet field can be produced at the LHC via the channels

$$\begin{aligned} q\bar{q} &\rightarrow Z^*/\Gamma^* \rightarrow \xi^{++}\xi^{--}, \\ q_1\bar{q}_2 &\rightarrow W^{\pm*} \rightarrow \xi^{++}\xi^{\mp}, \\ q\bar{q} &\rightarrow Z^*/\Gamma^* \rightarrow \xi^{+}\xi^{-}. \end{aligned} \quad (40)$$

In the above three channels the interactions are fixed by the triplet gauge couplings and hence the production cross sections only depend on the scalar masses. In addition to the above three channels, there are additional channels where the scalar triplet field can be produced in association with W^\pm or quarks,

$$\begin{aligned} q_1\bar{q}_2 &\rightarrow W^{\pm*} \rightarrow \xi^{++}W^\mp, \\ q_1q_2 &\rightarrow W^{\pm*}W^{\pm*}q_3q_4 \rightarrow \xi^{\pm\pm}q_3q_4, \\ q_1q_2 &\rightarrow Z^*Z^*q_3q_4 \rightarrow \xi^{\pm\pm}q_3q_4, \\ q_1q_2 &\rightarrow \Gamma^*\Gamma^*q_3q_4 \rightarrow \xi^{\pm\pm}q_3q_4. \end{aligned} \quad (41)$$

The associated production with W^\pm and single production via $W^\pm W^\pm$ fusion involve the $\xi^{\pm\pm}W^\pm W^\pm$ vertex, which is suppressed by a factor $\varepsilon = v_3/v_1$. The $\Gamma\Gamma$ and ZZ fusion processes are also very suppressed compared to the pair production cross section. The possible $\xi^{\pm\pm}$ decay modes are

$$\begin{aligned} \xi^{\pm\pm} &\rightarrow l_i^\pm l_j^\pm, \\ \xi^{\pm\pm} &\rightarrow W^\pm W^\pm, \\ \xi^{\pm\pm} &\rightarrow \xi^\pm W^\pm, \\ \xi^{\pm\pm} &\rightarrow \xi^\pm \xi^\pm, \end{aligned} \quad (42)$$

where $l_i = e, \mu, \tau$ for $i = 1, 2, 3$. The decay mode into pair of leptons have been extensively discussed in the literature because it provides a clear multi-lepton final state signatures for the pair production of doubly charged Higgs field with

a very small SM background [45]. The possible two body decay modes of ξ^\pm are

$$\begin{aligned} \xi^\pm &\rightarrow l_i^\pm \nu_j, \\ \xi^\pm &\rightarrow W^\pm Z, \\ \xi^\pm &\rightarrow u_j \bar{d}_k, \bar{u}_j d_k, \end{aligned} \quad (43)$$

with the last two decay modes again suppressed by a factor $\varepsilon = v_3/v_1$. Thus the production of scalar triplet fields can give rise to several possible final states. The final states can be classified according to the number of charged leptons as (a) $l^+l^+l^-l^-X$, (b) $l^\pm l^\pm l^\mp X$, (c) $l^\pm l^\pm X$, (d) $l^+l^-j_\tau X$, (e) $l^\pm j_\tau j_\tau X$, where l corresponds to electrons or muons (not necessarily all with the same flavor), j_τ corresponds to a tau jet and X represents additional jets [45]. The unique signature of model A is the decay mode $\xi^{++} \rightarrow \eta^+ \eta^+$, if kinematically allowed.

Similarly, in model B the dominant partonic production channels of the charged and neutral components of the fermion triplet are given by

$$\begin{aligned} q\bar{q} &\rightarrow Z^*/\Gamma^* \rightarrow \Sigma^+ \Sigma^-, \\ q_1\bar{q}_2 &\rightarrow W^{\pm*} \rightarrow \Sigma^\pm \Sigma^0. \end{aligned} \quad (44)$$

The decay modes of Σ^\pm, Σ^0 are

$$\begin{aligned} \Sigma^\pm &\rightarrow l^\pm Z, \\ \Sigma^\pm &\rightarrow l^\pm \Phi, \\ \Sigma^\pm &\rightarrow \bar{\nu} W^+, \nu W^-, \\ \Sigma^0 &\rightarrow l^\pm W^\mp, \\ \Sigma^0 &\rightarrow \nu Z, \\ \Sigma^0 &\rightarrow \nu \Phi. \end{aligned} \quad (45)$$

Here the final states with different no of leptons can be classified as (a) six leptons, (b) five leptons, (c) $l^\pm l^\pm l^\pm l^\mp X$, (d) $l^+l^+l^-l^-X$, (e) $l^\pm l^\pm l^\pm X$, (f) $l^\pm l^\pm l^\mp X$, (g) l^+l^-X , (h) l^+l^-jjjjX and (i) $l^\pm jjjjX$ [45, 46]. The unique signatures of the type III seesaw such as six lepton and five lepton final states can be used to distinguish it from the type II seesaw scheme at the LHC.

5 Conclusions

In conclusion, we have studied the realization of mass varying neutrinos in an extension of the usual triplet Higgs model by including an extra Higgs doublet (η) and an extension of the SM with fermion triplet (Σ_R). We find that both scenarios can accommodate ν DE with a dynamical neutrino mass related to the acceleron field while satisfying the naturalness, in the former scenario with TeV scale triplet Higgs fields (ξ) and additional doublet Higgs field (η) and in the latter scenario with TeV scale fermion triplets Σ . We also point out the possible leptogenesis mechanisms for simultaneously

generating the observed baryon asymmetry of the universe in both scenarios. Finally, the TeV scale new fields in both models give unique and highly predictive collider signatures, testable in the current run of LHC.

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