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Scalar triplet leptogenesis in the presence of right-handed neutrinos with S_3 symmetry

Subhasmita Mishra and Anjan Giri[1]

Department of Physics, IIT Hyderabad, Kandi-502285, India

E-mail: giria@iith.ac.in

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Abstract

Leptogenesis appears to be a viable alternative to account for the baryon asymmetry of the Universe through baryogenesis. In this context, we consider a scenario in which the standard model is extended with S_3 and Z_2 symmetry in addition to the two scalar triplets, two scalar doublets and three right-handed neutrinos. Presence of scalar triplets and right-handed neutrinos in the scenarios of both type-I and type-II seesaw frameworks provide a different leptogenesis option and can help us to understand the matter-antimatter asymmetry with simple S_3 symmetry. We discuss the neutrino phenomenology and leptogenesis in both high ($O(10^{10})$ GeV) and low energy scale (O(2) TeV) by constraining the Yukawa couplings. Moreover, we also consider the constraints on model parameters from neutrino oscillation data and leptogenesis to explain the rare lepton flavor violating decay and muon g-2 anomaly.

Keywords: neutrino mass, leptogenesis, scalar triplets

(Some figures may appear in colour only in the online journal)

1. Introduction

The Standard Model (SM) of particle physics has attained an unprecedented level of success over the last few decades, which culminated with the discovery of the Higgs boson at the CERN Large Hadron Collider. However, there appear to be observations which cannot be explained within the framework of the SM. In this context, the observation of neutrino oscillation has indicated that the SM needs to be extended to accommodate the massive neutrinos [1]. Moreover, there exists evidence of the baryon asymmetry of the Universe, with the obtained value of $\Omega_B h^2 = 0.0223 \pm 0.0002$ [2] that corresponds to the baryon asymmetry $Y_B \equiv \eta_B/s \approx 0.86 \times 10^{-10}$. There have been many attempts to find some hint of the physics

[1] Author to whom any correspondence should be addressed.

beyond the SM (BSM), but the quest so far remains unsuccessful. In the absence of any clear cut idea so as to ascertain the nature of the new physics, it is quite natural to explore simple extensions of the SM, which can help to explain the observed data.

The leptonic sector, in particular, the study of neutrinos has taken the center stage in particle physics in recent years. Discrete symmetries are widely used for a long time for BSM model building and to explain the neutrino phenomenology [3–9]. The discrete symmetries commonly discussed are the S_3, S_4 and A_4 symmetries to explain the observed neutrino oscillation data. Here we choose the simplest permutation symmetry, the S_3 along with Z_2 symmetry, to explain neutrino mass and also discuss leptogenesis [10–13]. In addition to the SM particle spectrum, we introduce three right-handed neutrinos, two Higgs doublets and two scalar triplets to explain the neutrino mass with type I+II seesaw mechanism [14]. There are a lot of studies using S_3 symmetry but here we would like to add another aspect of it to the growing list of possibilities. Earlier, it has been discussed in the literature that S_3 symmetry with type-I seesaw scenario could be helpful to accommodate the experimental findings in both quark and lepton sectors, in addition to explaining leptogenesis. Despite the simplicity of the type I seesaw model, type II seesaw is equally frequented due to the fact that addition of scalar does not lead to any anomaly, neither does it have negative contribution to the radiative correction of the SM Higgs mass, unlike the fermions. Leptogenesis with S_3 symmetry and right-handed neutrinos have been considered before in [15] and here we consider the possible effect of scalar triplets with S_3 symmetry.

The CP violation prescribed by the Kobayashi–Maskawa mechanism of the SM is not capable of explaining the observed matter antimatter asymmetry of the Universe and, therefore, lepton asymmetry plays a significant role here. Leptogenesis appears to be an elegant mechanism where the asymmetry generated in the leptonic sector of the SM can be converted to baryogenesis through sphaleron transitions and, in fact, this idea looks to be very promising. In general, lepton asymmetry produced by the out of equilibrium decay of right-handed neutrinos has been widely studied in the literature [16–36]. But, there are very few studies devoted to the generation of lepton asymmetry through the out of equilibrium decay of the scalar triplets in type II seesaw framework [37–42]. Nonzero CP asymmetry cannot be generated with one loop contribution in the presence of only one scalar triplet [38]. Hence the scalar sector should be extended with at least one more triplet to generate a nonzero CP asymmetry from the interference of tree and one loop contributions. Since the scalar triplet has two different decay modes, even though the gauge interactions and the total decay rate of the triplets are larger than the expansion rate of the Universe, still the lepton asymmetry can be generated with any of the decays being out of equilibrium. There exist studies in the literature in connection with the leptogenesis from scalar triplet in the presence of right-handed neutrinos [14, 43], and we focus here in this direction with some additional symmetries.

Motivated by the need to look for scenarios BSM, we chose the simplest discrete symmetry (S_3) with minimal particle contents to discuss neutrino phenomenology and leptogenesis as a viable option. Generation of lepton asymmetry from the decay of heavy Majorana fermions is well described in the literature in the framework of S_3 symmetry. But few studies have done with scalar triplets in this context. Therefore, we tried to explain the leptogenesis phenomena from the decay of heavy scalar triplet and its interesting phenomenology in the presence of right-handed Majorana neutrinos in a composite seesaw (type I+II) scenario.

The remainder of the article is as follows: in section 2, we explain the model framework. Here we describe the Lagrangian with S_3 symmetry, including the extended particle contents. Extended Higgs sector is also discussed in this section. Section 3 includes the description of the neutrino masses and mixing with type I+II seesaw mechanism, where we discuss the relevant framework to compare our results with the observed data. Section 4 is devoted to the

Table 1. Particle contents and quantum numbers under SM, S_3, and Z_2.

Particles	$SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$	S_3	Z_2
L_e, L_μ	(1, 2, −1)	2	+1
L_τ	(1, 2, −1)	1	+1
E_{1R}, E_{2R}	(1, 1 , −2)	2	+1
E_{3R}	(1, 1, −2)	1	+1
N_{1R}, N_{2R}	(1, 1, 0)	2	+1
N_{3R}	(1, 1, 0)	1	−1
H_1, H_2	(0, 2, 1)	2	+1
H_3	(0, 2, 1)	1	−1
Δ_1	(0, 3, 2)	1	+1
Δ_2	(0, 3, 2)	1	+1

Table 2. Neutrino oscillation parameters in global fit 3σ observation [53].

Parameter	3σ range
Δm^2_{21} $[10^{-5}\mathrm{eV}^2]$	6.79–8.01
$\lvert\Delta m^2_{31}\rvert$ $[10^{-3}\mathrm{eV}^2]$ (NO)	2.43–2.62
$\lvert\Delta m^2_{31}\rvert$ $[10^{-3}\mathrm{eV}^2]$ (IO)	2.41–2.60
$\sin^2\theta_{12}/10^{-1}$	2.75–3.5
$\sin^2\theta_{23}/10^{-1}$ (NO)	4.28–6.24
$\sin^2\theta_{23}/10^{-1}$ (IO)	4.33–6.23
$\sin^2\theta_{13}/10^{-2}$ (NO)	2.04–2.4
$\sin^2\theta_{13}/10^{-2}$ (IO)	2.04–2.46
$\delta_{\mathrm{CP}}(^\circ)$	125–392

leptogenesis. Here we discuss all the possible scenarios and associated Boltzmann equation including the results. Section 5 contains the conclusion.

2. The model

In this section, we discuss the particle content and corresponding group charges of the SM and extra particles, excluding the quark sector and focus only on the leptonic sector. The importance of discrete symmetries in particle phenomenology has already been discussed extensively in various studies earlier [4, 11]. We consider the extension of the SM ($SU(3) \times SU(2)_L \times U(1)_Y$)) with the simplest non-Abelian discrete flavor symmetry, S_3, and the Abelian symmetry Z_2. In addition to the SM particles, we include three right-handed neutrinos($N_{(1,2,3)R}$), two Higgs doublets, and two Higgs triplets ($\Delta_{1,2}$) to explain the neutrino mixing and leptogenesis.

In table 1 L_e, L_μ and L_τ are the first, second and third generation lepton families, respectively, N_{iR} and $\Delta_{1,2}$ are the right-handed singlet Majorana neutrinos and SU(2) triplet Higgs, respectively. The scalar triplets are defined in SU(2) basis and are given by [37]

$$\Delta_i = \begin{pmatrix} \frac{\Delta_i^+}{\sqrt{2}} & \Delta_i^{++} \\ \Delta_i^0 & -\frac{\Delta_i^+}{\sqrt{2}} \end{pmatrix}. \tag{1}$$

The invariant Lagrangian for both type I and type II, involving the scalars and fermions in the framework under consideration ($\mathrm{SU}(3)_c \otimes \mathrm{SU}(2)_L \otimes \mathrm{U}(1)_Y \otimes S_3 \otimes Z_2$), is given by [15, 44]

$$\begin{aligned}
\mathcal{L} \supset\ & y_{t1}[\overline{\tilde{L}_e}\Delta_1 L_e + \overline{\tilde{L}_\mu}\Delta_1 L_\mu] + y_{t1}'[\overline{\tilde{L}_\tau}\Delta_1 L_\tau] \\
& + y_{t2}[\overline{\tilde{L}_e}\Delta_2 L_e + \overline{\tilde{L}_\mu}\Delta_2 L_\mu] + y_{t2}'[\overline{\tilde{L}_\tau}\Delta_2 L_\tau] \\
& - y_{\nu_1}[\overline{L_e}\tilde{H}_2 N_{1R} + \overline{L_\mu}\tilde{H}_1 N_{1R} + \overline{L_e}\tilde{H}_1 N_{2R} - \overline{L_\mu}\tilde{H}_2 N_{2R}] \\
& - y_{\nu_3}[\overline{L_\tau}\,\tilde{H}_1 N_{1R} + \overline{L_\tau}\,\tilde{H}_2 N_{2R}] - y_{\nu_4}[\overline{L_\tau}\,\tilde{H}_3 N_{3R}] \\
& - y_{l2}[\overline{L_e}H_2 E_{1R} + \overline{L_\mu}H_1 E_{1R} + \overline{L_e}H_1 E_{2R} - \overline{L_\mu}H_2 E_{2R}] \\
& - y_{l4}[\overline{L_\tau}\,H_1 E_{1R} + \overline{L_\tau}\,H_2 E_{2R}] - y_{l5}[\overline{L_e}\,H_1 E_{3R} + \overline{L_\mu}\,H_2 E_{3R}] \\
& - \frac{1}{2}\textstyle\sum_{i=1,2}\overline{N}^c_{iR}M_{iR}N_{iR} - \frac{1}{2}\overline{N}^c_{3R}M_{3R}N_{3R} + \ \mathrm{h.c} - V(H_i, \Delta_j)\ \ (i = 1, 2, 3; j = 1, 2).
\end{aligned} \tag{2}$$

In the above expression, $\overline{\tilde{L}_L} = \overline{L^c{}_L}i\tau_2 = (-\overline{e^c{}_L} \quad \overline{\nu^c{}_L})$, y_{ν_i} and y_{li} are the Yukawa couplings of neutral and charged leptons, respectively. M_{iR} are the Majorana masses of right-handed neutrinos. Models with extra Higgs in the presence of discrete symmetries are well studied in the literature [45–52]. With the additional scalar content in the model, we can write the interaction potential as

$$\begin{aligned}
V(H_i, \Delta_j) =\ & m_0^2 H_3^\dagger H_3 + m_d^2(H_2^\dagger H_2 + H_1^\dagger H_1) + \lambda_1(H_2^\dagger H_2 + H_1^\dagger H_1)^2 \\
& + \lambda_2(H_1^\dagger H_2 - H_2^\dagger H_1)^2 + \lambda_3[(H_1^\dagger H_1 - H_2^\dagger H_2)^2 + (H_1^\dagger H_2 + H_2^\dagger H_1)^2] \\
& + \lambda_4[(H_3^\dagger H_1)(H_1^\dagger H_2 + H_2^\dagger H_1) + (H_3^\dagger H_2)(H_1^\dagger H_1 - H_2^\dagger H_2) + \mathrm{h.c}] \\
& + \lambda_5[(H_3^\dagger H_3)(H_1^\dagger H_1 + H_2^\dagger H_2)] + \lambda_6[(H_3^\dagger H_1)(H_1^\dagger H_3) + (H_3^\dagger H_2)(H_2^\dagger H_3)] \\
& + \lambda_7[(H_3^\dagger H_1)(H_3^\dagger H_1) + (H_3^\dagger H_2)(H_3^\dagger H_2) + \mathrm{h.c}] + \lambda_8(H_3^\dagger H_3)^2 \\
& + \mu_{SB1}^2(H_1^\dagger H_2 + \mathrm{h.c}) + \mu_{SB2}^2(H_3^\dagger(H_1 + H_2)) \\
& + m_{t1}^2\,\mathrm{Tr}(\Delta_1^\dagger\Delta_1) + m_{t2}^2\,\mathrm{Tr}(\Delta_2^\dagger\Delta_2) + \mu_1(\tilde{H}_1^\dagger\Delta_1^\dagger H_1 + \tilde{H}_2^\dagger\Delta_1^\dagger H_2) \\
& + \mu_1'(\tilde{H}_3^\dagger\Delta_1^\dagger H_3) + \mu_2(\tilde{H}_1^\dagger\Delta_2^\dagger H_1) + \mu_2'(\tilde{H}_3^\dagger\Delta_2^\dagger H_3) + g_2(H_3^\dagger\Delta_1^\dagger\Delta_1 H_3) \\
& + g_1(H_1^\dagger\Delta_1^\dagger\Delta_1 H_1 + H_2^\dagger\Delta_1^\dagger\Delta_1 H_2) + g_3(H_1^\dagger\Delta_2^\dagger\Delta_2 H_1) + g_4(H_3^\dagger\Delta_2^\dagger\Delta_2 H_3) \\
& + k_1((H_1^\dagger H_1)\mathrm{Tr}(\Delta_1^\dagger\Delta_1) + (H_2^\dagger H_2)\mathrm{Tr}(\Delta_1^\dagger\Delta_1)) + k_2((H_3^\dagger H_3)\mathrm{Tr}(\Delta_1^\dagger\Delta_1)) \\
& + k_3((H_1^\dagger H_1)\mathrm{Tr}(\Delta_2^\dagger\Delta_2) + (H_2^\dagger H_2)\mathrm{Tr}(\Delta_2^\dagger\Delta_2)) + k_4((H_3^\dagger H_3)\mathrm{Tr}(\Delta_2^\dagger\Delta_2)) \\
& + t_1\,\mathrm{Tr}\,(\Delta_1^\dagger\Delta_1)^2 + t_2\,\mathrm{Tr}\,(\Delta_2^\dagger\Delta_2)^2 + t_3\,\mathrm{Tr}(\Delta_1^\dagger\Delta_1)\mathrm{Tr}(\Delta_2^\dagger\Delta_2).
\end{aligned} \tag{3}$$

The minimization conditions are given by $\frac{\partial V}{\partial v_1} = 0$, $\frac{\partial V}{\partial v_3} = 0$, $\frac{\partial V}{\partial u_1} = 0$, $\frac{\partial V}{\partial u_2} = 0$, where, $\langle\Delta_{1,2}\rangle = \begin{pmatrix} 0 & 0 \\ u_{1,2} & 0 \end{pmatrix}$, $\langle H_1\rangle = \begin{pmatrix} 0 \\ v_1 \end{pmatrix}$, $\langle H_2\rangle = \begin{pmatrix} 0 \\ v_2 \end{pmatrix}$ and $\langle H_3\rangle = \begin{pmatrix} 0 \\ v_3 \end{pmatrix}$.

We found the stability conditions of the scalar potential by using the co-positivity criteria [54], which are given below

Table 3. Some sample benchmark points (BP) for the couplings are provided by using equations (48) and (47), which satisfy both neutrino mass and leptogenesis simultaneously.

Parameters	M_{Δ_1} (GeV)	M_N (GeV)	$\tilde{y}_{\nu_3}$	$y_{t1}{}'\mu_L$	$\sum m_\nu$ (eV)	$\epsilon_{\rm CP} = \epsilon_\Delta^N + \epsilon_\Delta^h$
BP1	10^{10}	2.1×10^{11}	0.02	1.9×10^3	0.07	7.6×10^{-7}
BP2	10^{10}	2.1×10^{11}	0.03	6.4×10^3	0.09	5.8×10^{-7}

$$
\begin{aligned}
&\lambda_1 + \lambda_3 \geqslant 0, \quad \lambda_8 \geqslant 0,\, l_1 \geqslant 0,\, l_2 \geqslant 0,\\
&\lambda_5 + \lambda_6 + |\lambda_7| + \sqrt{\lambda_8(\lambda_3 + \lambda_1)} \geqslant 0,\\
&3(\lambda_1 + \lambda_3)\sqrt{\lambda_8} + 2(\lambda_5 + \lambda_6 + |\lambda_7|)\sqrt{\lambda_1 + \lambda_3} \geqslant 0,\\
&2(\lambda_5 + \lambda_6 + |\lambda_7|)^2 - 3\lambda_8(\lambda_1 + \lambda_3) \geqslant 0,\\
&(g_1 + k_1) + \sqrt{(\lambda_1 + \lambda_3)t_1} \geqslant 0, \qquad (g_3 + k_3) + \sqrt{(\lambda_1 + \lambda_3)t_2} \geqslant 0,\\
&(g_2 + k_2) + \sqrt{\lambda_8 t_1} \geqslant 0, \quad (g_4 + k_4) + \sqrt{\lambda_8 t_2} \geqslant 0, \quad t_3 + \sqrt{t_2 t_1} \geqslant 0.
\end{aligned} \tag{4}
$$

The explicit symmetry breaking terms μ_{SB1} and μ_{SB2} in the potential in equation (3) break $S_3 \times Z_2$ softly to another symmetry, $H_1 \leftrightarrow H_2$. Hence one can choose, $\langle H_1^0 \rangle = \langle H_2^0 \rangle = v_1$. The symmetry breaking terms are important for generating masses of additional Higgs particles. Multi Higgs doublet models allow tree level flavor changing neutral current (FCNC), unless the coupling of scalar doublets to both up and down type quarks and leptons are protected. These FCNCs can be suppressed with Higgs masses at TeV scale, which cannot be generated by electroweak symmetry breaking. Therefore, one can adjust the heavy Higgs mass by fine-tuning the soft breaking parameters.

2.1. Masses and mixing in the Higgs sector

Looking at the symmetry breaking terms in the scalar potential in equation (3) one can redefine H_1 and H_2 in terms of H_+ and H_- as [15]

$$
H_1 = \frac{H_+ + H_-}{\sqrt{2}}, \quad H_2 = \frac{H_+ - H_-}{\sqrt{2}}. \tag{5}
$$

After redefinition, we can have $\langle H_-^0 \rangle = 0$ and $\langle H_+^0 \rangle = v_+$, with the assumption that $v_1 = v_2$. The mixing between H_+ and H_3 can be considered as both of them acquire a non zero vacuum expectation value (VEV). Here we can write the mass basis of these two Higgs fields by orthogonal rotation of flavor states as follows

$$
\begin{pmatrix} H_3 \\ H_+ \end{pmatrix} = \begin{pmatrix} \cos\beta & \sin\beta \\ -\sin\beta & \cos\beta \end{pmatrix} \begin{pmatrix} H_L \\ H_H \end{pmatrix}, \tag{6}
$$

where, $H_L = H_3 \cos\beta - H_+ \sin\beta$ and $H_H = H_3 \sin\beta + H_+ \cos\beta$, and β is the Higgs mixing angle. The new Higgs fields are written in SU(2) doublet form as

$$
H_L = \begin{pmatrix} 0 \\ h_L^0 + v \end{pmatrix}, \qquad H_H = \begin{pmatrix} h_H^+ \\ h_H^0 + i a_H \end{pmatrix}, \qquad H_- = \begin{pmatrix} h_-^+ \\ h_-^0 + i a_- \end{pmatrix}, \tag{7}
$$

where, H_L is the SM like Higgs with VEV, $v = \sqrt{v_+{}^2 + v_3{}^2} = 246$ GeV and $\tan\beta = \frac{v_+}{v_3}$ with $\langle H_3^0 \rangle = v_3$. Charged and CP odd components of H_L will be absorbed by the SM gauge bosons to acquire mass in unitary gauge conditions. And rest of the Higgs doublets will have two CP odd, two charged and two neutral scalar fields.

Table 4. Benchmark points for the parameters satisfying the constraints from neutrino mass and observed baryon asymmetry in TeV scale.

Parameters	M_{Δ_1}(TeV)	M_{Δ_2}(TeV)	$y_{t1}'\mu_{1L}$ (GeV)	$y_{t2}'\mu_{2L}$ (GeV)	$\sum m_\nu$(eV)	ϵ_{CP}
BP1	2	20	7.2×10^{-10}	9×10^{-7}	0.023	0.06
BP1	2	20	6×10^{-10}	7.6×10^{-7}	0.02	0.1

The masses of the scalar fields are given by

$$M_{h_H}^2 \approx M_{h_H^+}^2 \approx M_{a_H}^2 \approx \mu_{SB1}^2 \cos^2\beta + 2\sqrt{2}\mu_{SB2}^2 \sin\beta\cos\beta + (\mu_1 u_1 + \mu_2 u_2),$$
$$M_{h_L}^2 \approx \mathcal{O}(v^2),$$
$$M_{h_-}^2 \approx M_{h_-^+}^2 \approx M_{a_-}^2 \approx \mu_{SB1}^2 + \mu_1 u_1 + \mu_2 u_2.$$

The masses of the Higgs fields, other than the SM Higgs (125 GeV), can be achieved to be order of TeV by finetuning, which help in suppressing the tree level FCNCs.

2.1.1. Phase re-absorption. The phases of complex fermion fields can be redefined by fixing the phases in the complex Yukawa coupling present in the Lagrangian in equation (2) [15]. Let us consider that the neutral lepton Yukawa couplings transform as $y_{\nu_i} \to e^{ip_{yi}} y_{\nu_i}$ ($i = 1, 3, 4$), where, p_{yi} are the phases of transformations. Similarly for the charged lepton Yukawa and the fermion fields transform as $y_{li} \to e^{ip_{yli}} y_{li}$ ($i = 2, 4, 5$), $L_i \to e^{ip_l} L_i$ ($i = 1, 2$), $L_3 \to e^{ip_{l3}} L_3$, $E_{iR} \to e^{ip_E} E_{iR}$ ($i = 1, 2$), $E_{3R} \to e^{ip_{E3}} E_{3R}$, $N_{iR} \to e^{ip_R} N_{iR}$ ($i = 1, 2$) and $N_{3R} \to e^{ip_{3R}} N_{3R}$.

Phases of N_{iR} can be absorbed in the Majorana mass matrix M_{iR}, M_{3R} and the phases in the charged lepton Yukawa couplings can be fixed by the redefinition of the fermion fields, which can be found from the Lagrangian in equation (2) as follows

$$p_l = -p_{y_{l2}} + p_E, \quad p_{l3} = -p_{y_{l4}} + p_E, \quad p_{E3} = p_{y_{l5}} - p_{y_{l2}} + p_E. \tag{8}$$

Hence one can choose the leftover phase in charged lepton sector to be p_l, which can be fixed by the Yukawa coupling of the neutrinos, i.e. $p_l = -p_{y1}$. Therefore, the remnant phases in the neutral lepton complex Yukawa coupling are p_{y3} and p_{y4}. While constructing the mass matrices, the neutral and charged lepton fields can be rotated separately, hence only the relative phase $p_{y3} - p_{y4}$ appears in the neutrino mass matrix.

Similarly, in the scalar sector the triplet lepton Yukawa and the complex triplet fields transform as $y_{ti} \to e^{ip_{y_{ti}}} y_{ti}$ and $\Delta_i \to e^{ip_{\Delta_i}} \Delta_i$, respectively. From the triplet Lagrangian in equation (2) one can fix the phases as, $p_{y_{t1}} = 2p_l + p_{\Delta_1}$ and $p_{y_{t1}'} = 2p_{l3} + p_{\Delta_1}$. Hence, if the phases of y_{ti} can be absorbed by the redefinition of Δ_i, the remnant phases in the scalar triplet-lepton interaction sector are $p_{y_{ti'}}$.

3. Neutrino masses and mixing

In order to discuss the neutrino masses and mixing, we first discuss the type I seesaw mass matrix for neutral leptons, which is given in the basis of $\tilde{N} = (\nu_L^c, N_R)^T$ as

$$\mathcal{M} = \begin{pmatrix} 0 & M_D \\ M_D^T & M_R \end{pmatrix}. \tag{9}$$

We consider the light neutrino mass formula, which is described by the well known type I seesaw mechanism as [55, 56]

$$\mathcal{M}_\nu^{\mathcal{I}} = M_D M_R^{-1}(M_D)^{\mathrm{T}}. \tag{10}$$

Looking at the Lagrangian in equation (2), one can write the flavor structure of Dirac mass matrix for the neutral and charged leptons as

$$M_D = \begin{pmatrix} m_1 & m_1 & 0 \\ m_1 & -m_1 & 0 \\ m_3 & m_3 & m_4 \mathrm{e}^{\mathrm{i}\phi} \end{pmatrix}, \quad M_l = \begin{pmatrix} m_{l2} & m_{l2} & m_{l5} \\ m_{l2} & -m_{l2} & m_{l5} \\ m_{l4} & m_{l4} & 0 \end{pmatrix}, \tag{11}$$

where, $m_1 = y_{\nu_1} v_1$, $m_3 = y_{\nu_3} v_1$, and $m_4 = y_{\nu_4} v_3$. Similarly, $m_{l2} = y_{l2} v_1$, $m_{l5} = y_{l5} v_1$, and $m_{l4} = y_{l4} v_1$ with the assumption that $v_1 = v_2$. Using the mixing of the Higgs fields and the redefinition of VEVs, one can rewrite the terms inside the Dirac mass matrix as, $m_1 = y_{\nu_1} v \sin\beta$, $m_3 = y_{\nu_3} v \sin\beta$, and $m_4 = y_{\nu_4} v \cos\beta$. Similarly, one can also write, $m_{l2} = y_{l2} v \sin\beta$, $m_{l5} = y_{l5} v \sin\beta$, and $m_{l4} = y_{l4} v \sin\beta$. From the phase re-absorption, as explained before, we can always have the choice to put the remnant phase in any element of Dirac term. Here we put the relative phase ($\phi = p_{y3} - p_{y4}$) in m_4 for simplicity in calculation. Similarly, the type II mass matrix can be constructed from the Lagrangian in equation (2) by using the general formula of effective neutrino mass matrix in type II seesaw mechanism [57], i.e $\mathcal{M}_\nu^{\mathcal{II}} = \sum_i \frac{2\mu_{iL} y_{\Delta_i} v^2}{M_{\Delta_i}^2}$. The structure of mass matrix is given below

$$\mathcal{M}_\nu^{\mathcal{II}} = \begin{pmatrix} x_1 y_{t1} + x_2 y_{t2} & 0 & 0 \\ 0 & x_1 y_{t1} + x_2 y_{t2} & 0 \\ 0 & 0 & x_1 y_{t1}{}' + x_2 y_{t2}{}' \end{pmatrix}, \tag{12}$$

where, $x_i = \frac{2\mu_{iL} v^2}{M_{\Delta_i}^2}(i = 1, 2)$ and $\mu_{iL} = 2\mu_i \sin^2\beta + \mu_i{}' \cos^2\beta$. Here μ_{iL} is the coupling of triplets to the SM like Higgs(H_L) in this model, which contribute to the neutrino mass.

3.1. Diagonalization of charged lepton and neutrino mass matrices

The squared charged lepton mass matrix, which is Hermitian, can be diagonalized by unitary transformation as $U_{eL} M_l M_l{}^\dagger U_{eL}{}^\dagger = \mathrm{Diag}(m_e^2,\ m_\mu^2,\ m_\tau^2)$. The analytical diagonalization can be done by solving the characteristic polynomials of the given matrix, where the equations are provided in terms of squared charged lepton masses

$$\mathrm{Tr}(M_l M_l^\dagger) = m_e^2 + m_\mu^2 + m_\tau^2 = 4m_{l2}^2 + 2m_{l5}^2 + 2m_{l4}^2, \tag{13}$$

$$\det(M_l M_l^\dagger) = m_e^2 m_\mu^2 m_\tau^2 = 4m_{l2}^2 m_{l5}^2 m_{l4}^2, \tag{14}$$

$$\chi(M_l M_l^\dagger) = \frac{1}{2}[(\mathrm{Tr}[M_l M_l^\dagger])^2 - \mathrm{Tr}\,(M_l M_l^\dagger)^2] = 4[m_{l2}^4 + m_{l2}^2(m_{l4}^2 + m_{l5}^2) + m_{l4}^2 m_{l5}^2]. \tag{15}$$

Hence, we can get the solutions for m_{l2}, m_{l4} and m_{l5} in terms of charged lepton masses by solving the above equations

$$\frac{m_{l2}^2}{m_\tau^2} \approx \frac{1}{2}\frac{m_e^2 + m_\mu^2}{m_\tau^2} - \frac{m_e^2 m_\mu^2}{m_\tau^2(m_e^2 + m_\mu^2)} + p, \tag{16}$$

$$\frac{m_{l4,l5}^2}{m_\tau^2} \approx \frac{1}{4}\left(1 - \frac{m_e^2 + m_\mu^2}{m_\tau^2} + 4\frac{m_e^2 m_\mu^2}{m_\tau^2(m_e^2 + m_\mu^2)} - 4p\right) \pm \left[\left(1 - \frac{m_e^2 + m_\mu^2}{m_\tau^2} + 4\frac{m_e^2 m_\mu^2}{m_\tau^2(m_e^2 + m_\mu^2)} - 4p\right)^2 - \frac{m_\mu^2 m_e^2}{m_{l2}^2 m_\tau^2}\right]. \tag{17}$$

Here, p is the smallest root of the solution of the equation given in [58]. The eigenvector matrix of the squared masses can be obtained by solving the characteristic equation [11, 58] with an order approximation up to $\frac{m_e^2 m_\mu^2}{m_\tau^4}$

$$U_{el} = \begin{pmatrix} \frac{x}{\sqrt{2(1-x^2)}} & \frac{1}{\sqrt{2(1+x^2)}} & \frac{1}{\sqrt{2(1+\sqrt{z})}} \\ \frac{-x}{\sqrt{2(1-x^2)}} & \frac{-1}{\sqrt{2(1+x^2)}} & \frac{1}{\sqrt{2(1+\sqrt{z})}} \\ \frac{\sqrt{1-2x^2}}{\sqrt{1-x^2}} & \frac{x}{\sqrt{1+x^2}} & \frac{\sqrt{z}}{\sqrt{(1+\sqrt{z})}} \end{pmatrix}. \tag{18}$$

Where, $x = \frac{m_e}{m_\mu}$, and $z = \frac{m_e^2 m_\mu^2}{m_\tau^4}$. With the consideration of the Majorana neutrinos to be in diagonal basis, for simplicity, we can write the effective small neutrino mass matrix $\mathcal{M}_\nu^{\mathcal{I}}$ in type I seesaw framework from equations (10) and (11) as

$$\mathcal{M}_\nu^{\mathcal{I}} = \begin{pmatrix} \eta_1^2 & 0 & \eta_1\eta_3 \\ 0 & \eta_1^2 & 0 \\ \eta_1\eta_3 & 0 & \eta_3^2 + \eta_4^2 e^{2i\phi} \end{pmatrix}, \tag{19}$$

where, $\eta_1 = \frac{\sqrt{2}m_1}{\sqrt{M_{1R}}}$, $\eta_3 = \frac{\sqrt{2}m_3}{\sqrt{M_{1R}}}$, $\eta_4 = \frac{\sqrt{2}m_4}{\sqrt{M_{3R}}}$, and m_1, m_3 and m_4 are defined in equation (11). The small neutrino mass matrix for the type I+II seesaw scenario from equation (12) to (19) is given by $\mathcal{M}_\nu = \mathcal{M}_\nu^{\mathcal{I}} + \mathcal{M}_\nu^{\mathcal{II}}$, which can be written in matrix form as

$$\mathcal{M}_\nu = \begin{pmatrix} \eta_1^2 + y_{t1}x_1 + x_2 y_{t2} & 0 & \eta_1\eta_3 \\ 0 & \eta_1^2 + y_{t1}x_1 + x_2 y_{t2} & 0 \\ \eta_1\eta_3 & 0 & \eta_3^2 + \eta_4^2 e^{2i\phi} + |y_{t1}'x_1 + y_{t2}'x_2| e^{i\phi_\Delta} \end{pmatrix}. \tag{20}$$

With the consideration of y_{ti}', y_{ti}, μ_i to be complex, the phase in y_{ti} is fixed by the redefinition of field. The remnant phase in the triplet interaction sector being ϕ_Δ, which is the relative phase in y_{t1}' and y_{t2}'. Now the above mass matrix is reduced to a simple form as

$$\mathcal{M}_\nu = \begin{pmatrix} r_1 & 0 & \eta_1\eta_3 \\ 0 & r_1 & 0 \\ \eta_1\eta_3 & 0 & r_2 e^{i\phi_{eff}} \end{pmatrix}. \tag{21}$$

In the above, $r_1 = \eta_1^2 + x_1 y_{t1} + x_2 y_{t2}$ and $r_2 = |\eta_3^2 + \eta_4^2 e^{2i\phi} + |x_1 y_{t1}' + x_2 y_{t2}'| e^{i\phi_\Delta}|$ are the S_3 parameters and the effective phase (ϕ_{eff}) in the mass matrix, which is given by

$$\tan\phi_{eff} = \frac{\eta_4^2 \sin 2\phi + |x_1 y_{t1}' + x_2 y_{t2}'| \sin\phi_\Delta}{\eta_3^2 + \eta_4^2 \cos 2\phi + |x_1 y_{t1}' + x_2 y_{t2}'| \cos\phi_\Delta}. \tag{22}$$

Since the mass matrix is already in block diagonal form, it is easy to diagonalize the only non-diagonal block by simple orthogonal rotation. The rotation matrix in 3×3 dimensional form is given by

$$U_\nu = \begin{pmatrix} \cos\theta & 0 & \sin\theta e^{-i\phi_\nu} \\ 0 & 1 & 0 \\ -\sin\theta e^{i\phi_\nu} & 0 & \cos\theta \end{pmatrix}, \quad \text{and} \quad {U_\nu}^T \mathcal{M}_\nu U_\nu = \begin{pmatrix} m_{\nu_1} e^{i\phi_1} & 0 & 0 \\ 0 & m_{\nu_2} & 0 \\ 0 & 0 & m_{\nu_3} e^{i\phi_3} \end{pmatrix}. \tag{23}$$

After diagonalizing the neutrino mass matrix in equation (21), we can write the complex mass parameters of equation (23) as

$$\begin{aligned} M_{\nu_1} &= m_{\nu_1} e^{i\phi_1} = \frac{r_1 + r_2 e^{i\phi_{eff}}}{2} - \frac{1}{2}[(r_1 - r_2 e^{i\phi_{eff}})^2 + 4(\eta_1\eta_3)^2]^{1/2}, \\ M_{\nu_3} &= m_{\nu_1} e^{i\phi_3} = \frac{r_1 + r_2 e^{i\phi_{eff}}}{2} + \frac{1}{2}[(r_1 - r_2 e^{i\phi_{eff}})^2 + 4(\eta_1\eta_3)^2]^{1/2}, \\ M_{\nu_2} &= r_1, \quad \tan\phi_\nu = \frac{r_2 \sin\phi_{eff}}{[r_1 - r_2\cos\phi_{eff}]} \end{aligned} \tag{24}$$

where ϕ_ν is the only phase appearing in the neutrino mixing matrix in equation (23) and hence can be considered as Dirac like phase. ϕ_1 and ϕ_3 are the Majorana like phases in the mass matrix. Hence by the standard parameterization of PMNS matrix, which is $U_{\mathrm{PMNS}} = {U_{el}}^\dagger U_\nu$, we can construct the neutrino mixing matrix for this model from equation (18) to (23) as

$$U_{\mathrm{PMNS}} \approx \begin{pmatrix} \frac{x\cos\theta}{\sqrt{2(1-x^2)}} & \frac{1}{\sqrt{2(1+x^2)}} & \frac{e^{-i\phi_\nu} x \sin\theta}{\sqrt{2(1-x^2)}} \\ -\frac{x\cos\theta}{\sqrt{2(1-x^2)}} - \frac{e^{i\phi_\nu}\sin\theta}{\sqrt{2(1+\sqrt{z})}} & -\frac{1}{\sqrt{2(1+x^2)}} & \frac{\cos\theta}{\sqrt{2(1+\sqrt{z})}} - \frac{e^{-i\phi_\nu} x\sin\theta}{\sqrt{2(1-x^2)}} \\ \frac{\sqrt{1-2x^2}\cos\theta}{\sqrt{1-x^2}} - \frac{e^{i\phi_\nu}\sqrt{z}\sin\theta}{\sqrt{1+\sqrt{z}}} & \frac{x}{\sqrt{1+x^2}} & \frac{\sqrt{z}\cos\theta}{\sqrt{1+\sqrt{z}}} + \frac{e^{-i\phi_\nu}\sqrt{1-2x^2}\sin\theta}{\sqrt{1-x^2}} \end{pmatrix}. \tag{25}$$

The mixing angles can be obtained by comparing with the standard U_{PMNS} matrix as

$$\begin{aligned} \sin\theta_{13} &= |U_{e3}| \approx \frac{x\sin\theta}{\sqrt{2(1-x^2)}}, \quad \tan\theta_{12} = |\frac{U_{e2}}{U_{e1}}| \approx \frac{\sqrt{1-x^2}}{x\cos\theta\sqrt{1+x^2}}, \\ \tan\theta_{23} &= |\frac{U_{\mu3}}{U_{\tau3}}| \approx |\frac{\sqrt{2(1-x^2)}\cos\theta - \sqrt{2(1+\sqrt{z})}x\sin\theta e^{-i\phi_\nu}}{\sqrt{2z(1-x^2)}\cos\theta + \sqrt{(1-2x^2)}\sqrt{2(1+\sqrt{z})}\sin\theta e^{-i\phi_\nu}}|. \end{aligned} \tag{26}$$

For a sample value of $\theta = \frac{\pi}{6}$, we found $\sin\theta_{13} \approx 0.004$, $\sin^2\theta_{12} \approx 0.6$ and $\sin^2\theta_{23} \approx 0.42$.

3.2. Numerical analysis

Now redefining the parameters appearing in the complex neutrino masses in equation (24) as $\rho_1 = \frac{r_2}{r_1}$ and $\rho_2 = \frac{\eta_1\eta_3}{r_1}$. We have the physical masses of the active neutrinos and corresponding Majorana phases in terms of new parameters as

$$m_{\nu_1} = |M_{\nu_1}| = \left|\frac{r_1}{2}\right| [(1 + \rho_1\cos\phi_{eff} - C)^2 + (\rho_1\sin\phi_{eff} - D)^2]^{1/2}, \tag{27}$$

$$m_{\nu_3} = |M_{\nu_3}| = \left|\frac{r_1}{2}\right| [(1 + \rho_1\cos\phi_{eff} + C)^2 + (\rho_1\sin\phi_{eff} + D)^2]^{1/2}, \tag{28}$$

$$m_{\nu_2} = |M_{\nu_2}| = |r_1|, \tag{29}$$

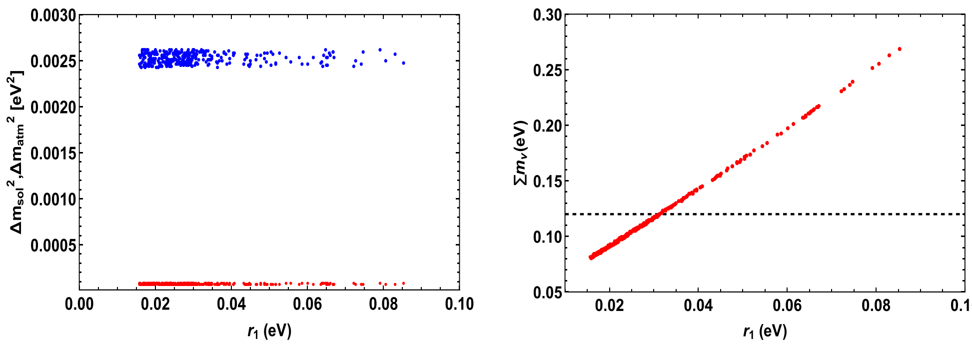

Figure 1. Variation of r_1 with the total active neutrino mass in the left panel and variation with solar (red) and atmospheric (blue) mass squared differences in the right panel.

$$\tan\phi_1 = \frac{\rho_1 \sin\phi_{eff} - D}{1 + \rho_1 \cos\phi_{eff} - C}, \quad \tan\phi_3 = \frac{\rho_1 \sin\phi_{eff} + D}{1 + \rho_1 \cos\phi_{eff} + C}, \tag{30}$$

where

$$C = \left(\frac{A + \sqrt{A^2 + B^2}}{2}\right)^{1/2}, \quad D = \left(\frac{-A + \sqrt{A^2 + B^2}}{2}\right)^{1/2},$$
$$A = 1 - 2\rho_1 \cos\phi_{eff} + \rho_1^2 \cos 2\phi_{eff} + 4\rho_2^2, \quad B = -2\rho_1 \sin\phi_{eff} + \rho_1^2 \sin 2\phi_{eff} \tag{31}$$

and ϕ_1 and ϕ_2 are the Majorana like phases. We prefer a normal ordering of active neutrino masses expressed in the above equations as functions of S_3 parameters

$$r = \left[\frac{\Delta m_{\rm sol}^2}{\Delta m_{\rm atm}^2}\right] = \frac{1}{4}\left(2 - \frac{-3 + C^2 + D^2 + {\rho_1}^2 + 2\rho_1 \cos\phi_{eff}}{C + C\rho_1 \cos\phi_{eff} + D\rho_1 \sin\phi_{eff}}\right) \approx 0.03. \tag{32}$$

We vary the model parameters as follows: r_1 from 0 to 0.1, ρ_1 from 0 to 1.5 and ρ_2 from 0 to 3, ϕ_{eff} between $-\pi$ to $+\pi$. We now show the allowed parameter space, compatible with the 3σ data of neutrino observables (table 2) and cosmological bound on total active neutrino mass. The left and right panel of figure 1 represents the variation of r_1 with solar and atmospheric mass squared differences, and total active neutrino mass, respectively. It gives a strong constraint on r_1 to lie within a range 0.015–0.03. The allowed parameter region for ρ_1, ρ_2 and ϕ_{eff} are represented in the left, middle and right panels of figure 2, respectively. We noticed that the favorable region for ρ_1 and ϕ_{eff} are 0 to 1.5 and -3.14 to $+3.14$, respectively, where the region below 1.3 is excluded for ρ_2. We predict the range on physical neutrino masses given in equations (27)–(29), shown in the left panel of figure 3. Here, we found the masses (in eV scale) $m_{\nu 1}$ to vary from 0.01 to 0.03, 0.015 to 0.03 for $m_{\nu 2}$ and 0.05 to 0.06 for $m_{\nu 3}$. The right panel shows the constraint on Majorana like phases, ϕ_1 to vary from -0.4 to 0.8 and ϕ_3 vary from -0.8 to 1.8 from neutrino oscillation data. In figure 4 we have shown the correlation between r_1 and ρ_2 (left panel), ϕ_ν with ϕ_{eff} (middle panel) and ϕ_ν with ρ_1 (right panel). However in figure 5, we represented the allowed region of Dirac like CP phase ϕ_ν,

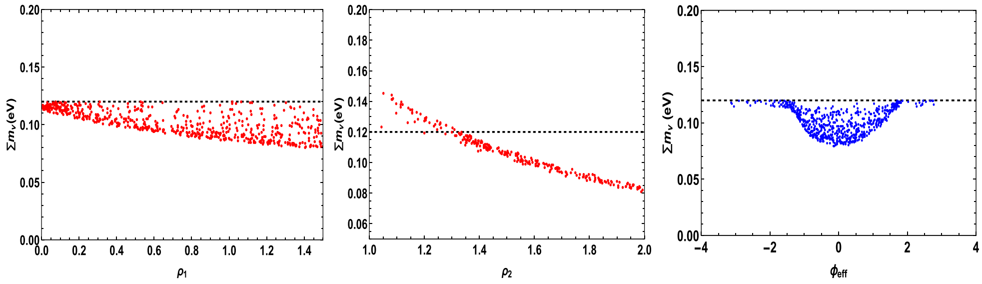

Figure 2. Variation of ρ_1 (left panel) and ρ_2 (middle panel) and ϕ_{eff} (right most panel) with the total active neutrino mass.

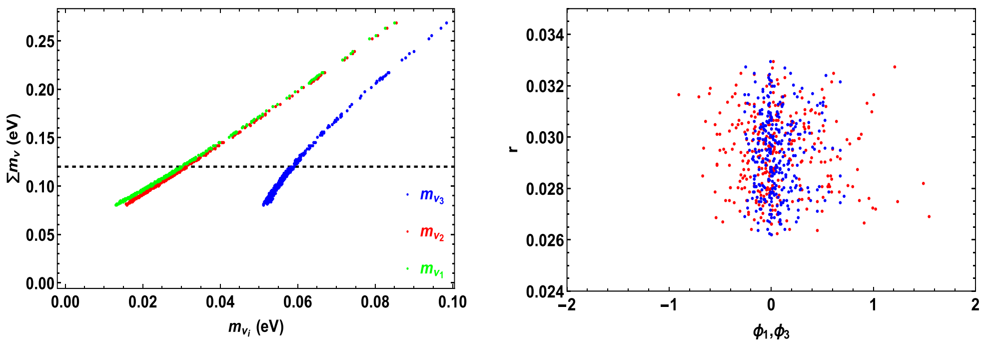

Figure 3. Left panel displays a variation of the physical neutrino masses with total active neutrino mass and variation of ϕ_1 (red) and ϕ_3 (blue) with ratio of solar to atmospheric mass squared difference are presented in the right panel.

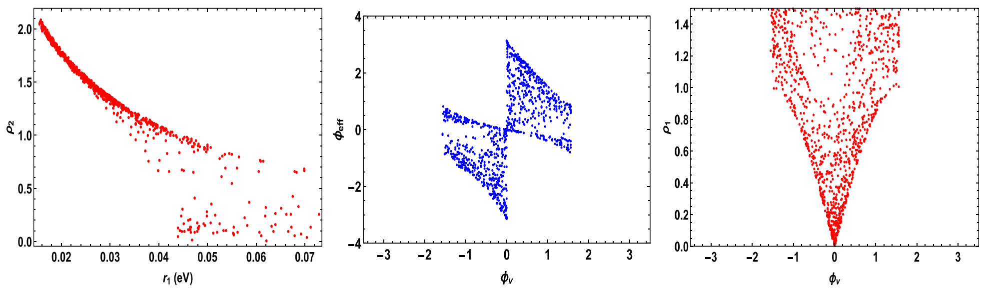

Figure 4. Left panel shows a correlation between r_1 and ρ_2, where the correlation of Dirac like phase (ϕ_ν) with ϕ_{eff} and ρ_1 are represented in the middle and right panel, respectively.

which lie in the region -1.5 to 1.5 in the left panel and in the right panel we get a more constrained value of ϕ_ν to vary from -0.59 to -1.5 compatible with the 2σ constraint on the CP phase from T2K data [59][2].

[2] While constraining the model parameters with consideration of 2σ range of CP violating phase from T2K, we found only ϕ_{eff} and ρ_1 are restricted to lie in the region -1.5 to 1 and 0.5 to 1.5, respectively, however the other parameters do not change appreciably.

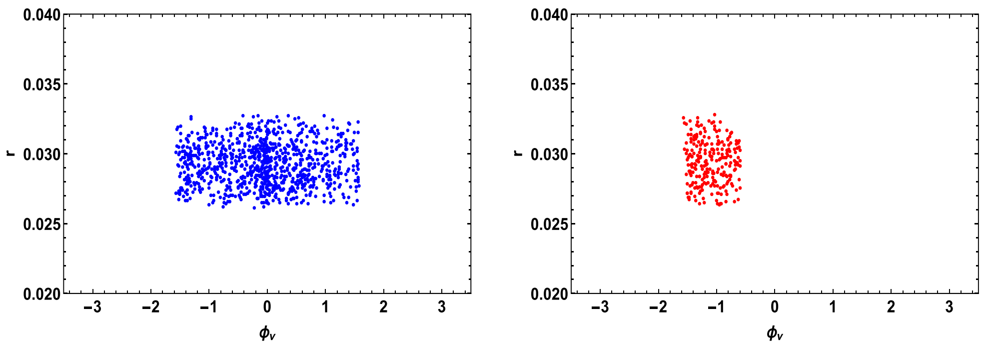

Figure 5. The left panel displays the prediction of Dirac CP phase (ϕ_ν) from global fit data and the right panel shows a more constrained parameter space for ϕ_ν with consideration of 2σ range of CP violating phase from T2K [59].

4. Leptogenesis

Leptogenesis, through the out of equilibrium decay of heavy particles, is the most viable way to generate observed baryon asymmetry of the Universe. For a review of leptogenesis, one can refer to [60]. Even though this formalism is widely studied with the heavy Majorana fermions, there exist few studies, which focus on the production of asymmetry through scalar decays within the type II seesaw framework [14, 37]. Here, as emphasized before, the scalar triplets are not enough to explain the observed neutrino mixing within only type II seesaw scenario. Therefore, we add extra right-handed neutrinos to include type I seesaw framework for the explanation of neutrino phenomenology. Unlike the Majorana fermions, triplets are charged and can generate an asymmetry in particle-antiparticle decay width. Both scalar triplets and right-handed neutrinos in the scenario under consideration can contribute to the leptonic CP violation. Various cases are possible, which include the generation of asymmetry solely from the decay of scalar triplets and right-handed neutrinos or from both, depending on their masses and are well studied in the literature [43]. In this work, we aim to explore the leptogenesis phenomenon from the lightest heavy triplet in a scenario of high and low energy scales. It is pointed out earlier in the literature that one cannot generate a non-zero CP asymmetry with one loop contribution in the presence of a single scalar triplet [38]. However, adding at least one additional triplet scalar can help resolve the problem. Before focusing on the one loop level decay of scalar triplets, we discuss the possible tree level decay channels, as the CP asymmetry is calculated from the interference of tree and the one loop level diagrams. From the type II seesaw Lagrangian equation (2), one can find two possible decay modes of the scalar triplets, namely, the triplet can decay to two leptons or two scalars in the final state, which are given by $\Delta \to \bar{\ell}_i \bar{\ell}_i$ and $\Delta \to h_j h_j$, where h_j are the Higgs fields of the model.

Summing over all the final states for leptons and Higgs decay modes of scalar triplet, we have

$$\Gamma(\Delta_i^{--} \to \ell^- \ell^-) = \Gamma(\Delta_i^{0\star} \to \nu\nu) = \Gamma(\Delta_i^- \to l^- \nu) = \frac{M_{\Delta i}}{8\pi} \mathrm{Tr}[\mathcal{Y}_{\Delta_i} \mathcal{Y}_{\Delta_i}^\dagger], \tag{33}$$

$$\Gamma(\Delta_i^{--} \to h^- h^-) = \Gamma(\Delta_i^{0\star} \to h_0{}^\star h_0{}^\star) = \Gamma(\Delta_i^- \to h^- h_0{}^\star) = \frac{1}{8\pi M_{\Delta_i}} |\mu_{\Delta_i}|^2. \tag{34}$$

Hence the branching ratios for the decay of Δ_1 are given by

$$B_l = \sum_{l=e,\mu,\tau} B_l = \frac{M_{\Delta_1}}{8\pi\Gamma_{tot}} Tr(\mathcal{Y}_{\Delta_1}\mathcal{Y}^{\dagger}_{\Delta_1}), \tag{35}$$

$$B_h = \sum_{i\ =\ L,H,-} B_{H_i} = \frac{|\mu_{1L}|^2 + |\mu_{1H}|^2 + |\mu_{1-}|^2}{8\pi M_{\Delta_1}\Gamma_{\Delta_1}} = \frac{|\lambda|^2 M_{\Delta_1}}{8\pi\Gamma_{\Delta_1}}, \tag{36}$$

$$\Gamma_{\Delta_1} = \frac{M_{\Delta_1}}{8\pi}\Big(\sum|\mathcal{Y}_{\Delta_1}|^2_{ii} + |\lambda|^2\Big), \tag{37}$$

where $\lambda = \frac{(|\mu_{1L}|^2 + |\mu_{1H}|^2 + |\mu_{1-}|^2)^{1/2}}{M_{\Delta_1}}$, and Γ_{Δ_1} is the total decay rate of the first triplet. μ_{1L}, μ_{1H} and μ_{1-} are the coupling of the decaying triplet to H_L, H_H and H_-, respectively, which are defined as follows

$$\begin{aligned}\mu_{1L} &= \mu_1 \sin^2\beta + {\mu_1}' \cos^2\beta,\\ \mu_{1H} &= \mu_1 \cos^2\beta + {\mu_1}' \sin^2\beta,\\ \mu_{1-} &= \mu_1.\end{aligned} \tag{38}$$

From the type II Lagrangian mentioned in equation (2), the scalar triplet-lepton Yukawa couplings can be written in the matrix form as

$$\mathcal{Y}_{\Delta_1} = \begin{pmatrix} y_{t1} & 0 & 0\\ 0 & y_{t1} & 0\\ 0 & 0 & {y_{t1}}' \end{pmatrix}, \qquad \mathcal{Y}_{\Delta_2} = \begin{pmatrix} y_{t2} & 0 & 0\\ 0 & y_{t2} & 0\\ 0 & 0 & {y_{t2}}' \end{pmatrix}. \tag{39}$$

Since the SM leptons and right-handed neutrinos couple to different Higgs doublets of the present model, we can re-write the Yukawa coupling matrix corresponding to the interaction of these fermions with H_L, H_H and H_- from equation (2) after redefinition and rotation of Higgs fields in equations (5) and (6). One can write

$$\mathcal{Y}^{\nu}_L = \begin{pmatrix} \frac{y_{\nu2}}{\sqrt{2}}\sin\beta & \frac{y_{\nu2}}{\sqrt{2}}\sin\beta & 0\\ \frac{y_{\nu2}}{\sqrt{2}}\sin\beta & -\frac{y_{\nu2}}{\sqrt{2}}\sin\beta & 0\\ \frac{y_{\nu3}e^{ip_{y3}}}{\sqrt{2}}\sin\beta & \frac{y_{\nu3}e^{ip_{y3}}}{\sqrt{2}}\sin\beta & y_{\nu_4}e^{ip_{y4}}\cos\beta \end{pmatrix}, \quad \mathcal{Y}^{\nu}_- = \begin{pmatrix} -\frac{y_{\nu2}}{\sqrt{2}} & \frac{y_{\nu2}}{\sqrt{2}} & 0\\ \frac{y_{\nu2}}{\sqrt{2}} & \frac{y_{\nu2}}{\sqrt{2}} & 0\\ \frac{y_{\nu3}}{\sqrt{2}} & -\frac{y_{\nu3}}{\sqrt{2}} & 0 \end{pmatrix},$$

$$\mathcal{Y}^{\nu}_H = \begin{pmatrix} \frac{y_{\nu2}}{\sqrt{2}}\cos\beta & \frac{y_{\nu2}}{\sqrt{2}}\cos\beta & 0\\ \frac{y_{\nu2}}{\sqrt{2}}\cos\beta & -\frac{y_{\nu2}}{\sqrt{2}}\cos\beta & 0\\ \frac{y_{\nu3}e^{ip_{y3}}}{\sqrt{2}}\cos\beta & \frac{y_{\nu3}e^{ip_{y3}}}{\sqrt{2}}\cos\beta & e^{ip_{y4}}y_{\nu_4}\sin\beta \end{pmatrix}. \tag{40}$$

Here we consider the leptogenesis in the mass basis of charged leptons. Hence the Dirac Yukawa matrices, mentioned above, will be modified as $\tilde{\mathcal{Y}}^{\nu}_{H,L,-} = \mathcal{Y}^{\nu}_{H,L,-}U_{el}$. If we use the observed masses of the charged leptons, we can have the numerical entries of the matrix U_{el} in equation (18), from which the modified Yukawa coupling for H_L as

$$\tilde{\mathcal{Y}}_L^\nu = \begin{pmatrix} 0 & 0 & y_{\nu_1}\sin\beta \\ 0.005y_{\nu_1}\sin\beta & y_{\nu_1}\sin\beta & 0 \\ y_{\nu_4}\cos\beta e^{ip_{y_4}} & 0.005y_{\nu_1}\sin\beta & 0.000\,01y_{\nu_4}\cos\beta e^{ip_{y_4}} + y_{\nu_3}\sin\beta e^{ip_{y_3}} \end{pmatrix}. \tag{41}$$

Similarly, we can have the modified Yukawa coupling matrices $\tilde{\mathcal{Y}}_H^\nu$ and $\tilde{\mathcal{Y}}_-^\nu$, corresponding to the Yukawa couplings of H_H and H_-. Further, the entries of modified Yukawa couplings will be denoted as $\tilde{y}_{\nu_i}$. We explore the importance of the above mentioned couplings in explaining leptogenesis in both high and low scale regimes in the following subsections.

4.1. Case I: Leptogenesis with $M_{\Delta_1} = \mathcal{O}(10^{10})$ GeV

Here we consider leptogenesis solely from scalar triplets by assuming the right-handed neutrinos to be much heavier in mass and decoupled earlier [43]. In the high temperature regime, the scalar triplets are thermalized because of the gauge interactions. They decouple, when the temperature of thermal bath approaches the mass of decaying triplet, i.e. $T \approx M_{\Delta_1}$ and produce the lepton asymmetry. We choose the mass of the scalar triplet to be of the order of $\mathcal{O}(10^{10})$ GeV and the mass of the lightest right-handed neutrino to be of the order of $\mathcal{O}(10^{11})$ GeV. Generation of nonzero CP asymmetry from the scalar triplet decay, in one loop level, requires at least one more heavy triplet ($M_{\Delta_2} \gg M_{\Delta_1}$), as mentioned in the literature. As the scalar triplet has two decay modes (to different Higgs in the model and the leptons), any of the decay channels being out of equilibrium can generate lepton asymmetry.

Constraints on couplings from out of equilibrium decay and neutrino mass. To generate nonzero asymmetry as per the Sakharov's condition, the decay rate should be less than the Hubble expansion

$$\begin{aligned}\Gamma_{\Delta_1} < H &\approx 1.66 \times \sqrt{g_\star}\frac{T^2}{1.2 \times 10^{19}} \text{ GeV} \\ &< 1.38 \times 10^{-18}[T^2]_{T=M_\Delta}(g_\star \approx 100) \text{ GeV}.\end{aligned} \tag{42}$$

The cosmological bound on sum of the neutrino masses is found to be less than 0.12 eV [61] . In this model, we are considering the total contribution to the neutrino mass contributing equally from type I and II sectors. Hence, we can constrain the Yukawa couplings with the assumption that the lightest right-handed neutrino is $\mathcal{O}(1)$ lighter in mass than the other two heavy neutrinos,

From type I:

$$\begin{aligned}\sum m_\nu^I &= \frac{4y_{\nu_1}^2 v^2\sin^2\beta}{M_{1R}} + \frac{2y_{\nu_3}^2 v^2\sin^2\beta}{M_{1R}} + \frac{2y_{\nu_4}^2 v^2\cos^2\beta}{M_{3R}} \leqslant 0.05 \times 10^{-9} \text{ GeV} \\ &\approx 0.3y_{\nu_3}^2\sin^2\beta + y_{\nu_4}^2\cos^2\beta \leqslant \frac{(0.05 \times 10^{-9})M_{3R}}{2v^2}.\end{aligned} \tag{43}$$

From type II:

$$\begin{aligned}\sum_i m_{\nu i}^{II} &= \frac{(4y_{ti} + 2y_{ti}{}')\mu_{iL}v^2}{M_{\Delta_i}^2} \leqslant 0.05 \times 10^{-9} \text{ GeV} \\ &\approx (2y_{ti} + y_{ti}{}')\mu_{iL} \leqslant \frac{(0.05 \times 10^{-9})M_{\Delta_i}^2}{2v^2} \text{ GeV}.\end{aligned} \tag{44}$$

The general expression for CP asymmetry from the leptonic self energy and vertex contribution to the scalar triplet decay is provided below

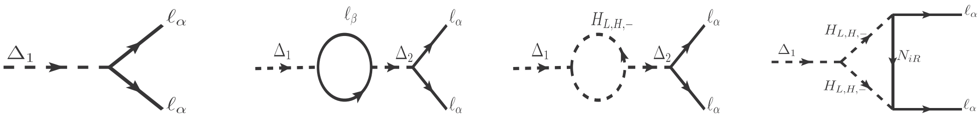

Figure 6. The scalar triplet tree and 1-loop diagrams in the presence of additional triple scalars contributing to the leptogenesis.

$$\epsilon_\Delta^\ell = \sum_i \frac{M_{\Delta_1}^2}{2\pi} \frac{\mathrm{Im}\,[(\mathcal{Y}_{\Delta_1}^\dagger \mathcal{Y}_{\Delta_2})_{ii}\,\mathrm{Tr}(\mathcal{Y}_{\Delta_1}^\dagger \mathcal{Y}_{\Delta_2})]}{M_{\Delta_1}^2\,\mathrm{Tr}[\mathcal{Y}_{\Delta_1}\mathcal{Y}_{\Delta_1}^\dagger] + \mu_{\Delta_1}^2} g\left(\frac{M_{\Delta_1}^2}{M_{\Delta_2}^2}\right), \tag{45}$$

where

$$g(x) = \frac{x(1-x)}{(1-x)^2 + xy}, \quad x = \frac{M_{\Delta_1}^2}{M_{\Delta_2}^2} \text{ and } y = \left(\frac{\Gamma_{\Delta_2}}{M_{\Delta_2}}\right)^2. \tag{46}$$

Hence the total CP asymmetry from the lepton loop in one flavor regime is $\epsilon_l = \sum_{\alpha=e,\mu,\tau}\epsilon_l^\alpha = 0$, due to the diagonal triplet-lepton Yukawa matrix. This structure of Yukawa also disfavors the flavored leptogenesis scenario. Hence the only self energy diagram that contributes to the triplet leptogenesis scenario is through Higgs loop, as shown in figure 6. Here the scalar triplet to Higgs decay mode to be out of equilibrium as per the Sakharov's conditions. This leads to $B_h\Gamma_{tot} < H$, where H is the Hubble expansion rate of the Universe. CP asymmetry from the interference of tree and scalar self energy loop is given by

$$\begin{aligned}
\epsilon_\Delta^h &= \textstyle\sum_i \frac{1}{2\pi} \frac{\mathrm{Im}\,(\mathcal{Y}_{\Delta_1}^\dagger \mathcal{Y}_{\Delta_2})_{ii}\mu_{\Delta_1}^\star \mu_{\Delta_2}}{M_{\Delta_1}^2 Tr[\mathcal{Y}_{\Delta_1}^\dagger \mathcal{Y}_{\Delta_1}] + |\mu_{\Delta_1}|^2} g\left(\frac{M_{\Delta_1}^2}{M_{\Delta_2}^2}\right) \\
&\approx \frac{1}{16\pi^2} \frac{\mathrm{Im}[y_{t1}{}' y_{t2}{}'(\mu_{1L}\mu_{2L}^\star + \mu_{1H}\mu_{2H}^\star + \mu_{1-}\mu_{2-}^\star)]}{\Gamma_{\Delta_1} M_{\Delta_1}} g\left(\frac{M_{\Delta_1}^2}{M_{\Delta_2}^2}\right) \\
&\approx \frac{3}{16\pi^2} \frac{3 y_{t1}{}' y_{t2}{}'(\mu_{1L}\mu_{2L}^\star)\sin\phi_{\mathrm{cph}}}{\Gamma_{\Delta_1} M_{\Delta_1}} g\left(\frac{M_{\Delta_1}^2}{M_{\Delta_2}^2}\right).
\end{aligned} \tag{47}$$

Since there is no contribution to CP asymmetry from the lepton loop in one flavor approximation, one can generate a large CP asymmetry from scalar self energy loop and vertex diagrams with right-handed neutrinos. CP asymmetry from the right-handed neutrino loop is given by

$$\begin{aligned}
\epsilon_\Delta^N &= \frac{-1}{4\pi} \textstyle\sum_i M_{iR} \frac{Im[\mu_{\Delta_1}(\mathcal{Y}_{\Delta_1})_{ii}(\tilde{\mathcal{Y}}_{L,H,-}^{\nu\star} \tilde{\mathcal{Y}}_{L,H,-}^{\nu\star})_{ii}]}{M_{\Delta_1}^2 Tr[\mathcal{Y}_{\Delta_1}\mathcal{Y}_{\Delta_1}^\dagger] + \mu_{\Delta_1}^2} \ln\left(1 + \frac{M_{\Delta_1}^2}{M_{iR}^2}\right) \\
&\approx \frac{-1}{4\pi} \frac{M_{\Delta_1}^2}{M_{3R}} \frac{\mu_{\Delta_1} y_{t1}{}' \tilde{y}_{\nu_3}{}^2 \sin\phi_{\mathrm{cpn}}}{M_{\Delta_1}{}^2 Tr(\mathcal{Y}_{\Delta_1}{}^\dagger \mathcal{Y}_{\Delta_1}) + |\mu_{\Delta_1}|^2} \quad (M_{iR} \gg M_{\Delta_1}).
\end{aligned} \tag{48}$$

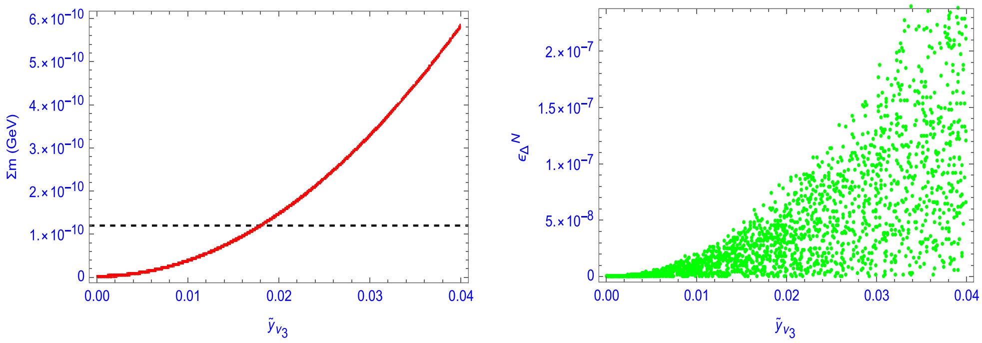

Figure 7. Left panel shows the variation of Yukawa coupling with the sum of neutrino masses in compatible with the 3σ neutrino oscillation parameters and cosmological bound on total neutrino mass. Right panel shows the variation of same Yukawa coupling with the CP asymmetry generated, satisfying the observed baryon asymmetry of the Universe.

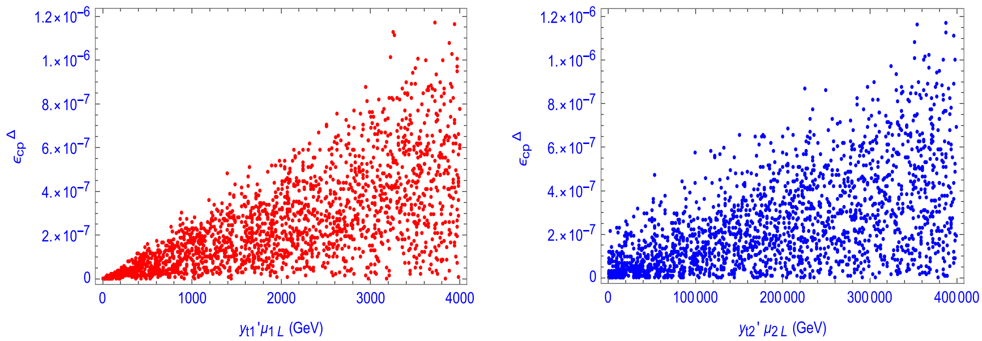

Figure 8. Left panel displays the variation of total CP asymmetry, generated by the first triplet with the corresponding triplet-lepton Yukawa coupling and right panel represents the variation of CP asymmetry with the Yukawa coupling for second triplet, compatible with the observed Baryon asymmetry and neutrino mass.

Here, we consider the constraint on the relevant Dirac Yukawa coupling (y_{ν_3}) from type I seesaw neutrino mass, by fixing the Higgs mixing angle $\beta = \frac{\pi}{4}$ and calculating the CP asymmetry by using equation (48). We show the variation of this coupling with the CP asymmetry in figure 7, from which one can infer that for low values of the Yukawa couplings the CP asymmetry turns out to be very small. But considering a range of 0.02–0.04, one can achieve a large CP asymmetry of the order 10^{-7}–10^{-8}, which is required for successful leptogenesis. Similarly, in figure 8, we have shown the variation of triplet-lepton Yukawa coupling, constrained from the type II neutrino mass and out of equilibrium decay width of the scalar triplet. A favored region of the coupling $y'_{ti}\mu_{iL}$ to vary within 1–3 TeV for the decaying triplet and 100–200 TeV for the heavier triplet scalar to obtain a CP asymmetry of order 10^{-7}. Some sample bench mark points for couplings those satisfy both neutrino mass and leptogenesis are provided in table 3.

4.1.1. Boltzmann equations. Efficiency of leptogenesis plays a vital role in generating the final baryon asymmetry, which could be governed by the dynamics of relevant Boltzmann equations. The particle dynamics in the early universe indulge a large number of interactions in the thermal soup. Particles attain thermal equilibrium and are subjected to the chemical equilibrium constraints, while the gauge interaction rate is more than the Hubble expansion. Here the chemical potential becomes important to define the relations between the particles in the chemical equilibrium. For the interactions that are equal with the Hubble expansion are not that fast to remain in equilibrium. Therefore the Boltzmann equations are very much significant to analyze the particle number density after the chemical or kinetic decoupling of particles in a specific temperature regime. In this model, we consider two Higgs triplets, where one of the Higgs triplets is more massive than the other. Hence the asymmetry generated by the heavier triplet will be washed out by the inverse decay of the lighter one. As it has been demonstrated in the literature earlier that the lepton number violation demands both the decay modes of the triplets to happen and any of the decay modes needs to be out of equilibrium to satisfy the Sakharov's condition. Unlike the right-handed neutrinos, as the scalar triplet is not Majorana particle and hence there will be asymmetry in particle and antiparticle decays and will contribute to the total asymmetry

$$H(T) = \frac{4\pi^3 g_\star}{45}\frac{T^2}{M_{pl}}, \quad \text{where,} \quad M_{pl} = 1.2 \times 10^{19}\,\text{GeV}, \tag{49}$$

$$\Sigma_\Delta^{eq} = \frac{45 g_T}{4\pi^4 g_\star} z^2 K_2(z), \quad Y_N^{eq} = \frac{135\zeta(3) g_N}{16\pi^4 g_\star} z^2 K_2(z), \tag{50}$$

$$Y_l^{eq} = \frac{3}{4}\frac{45\zeta(3) g_l}{2\pi^4 g_\star}, \quad Y_h^{eq} = \frac{45\zeta(3) g_h}{2\pi^4 g_\star}, \tag{51}$$

where, $g_\star = 106.75$ is the total relativistic degree of freedom of the SM particles in the equilibrium. $g_l = 2$, $g_h = 1$, $g_T = 1$, $g_N = 2$ are the degrees of freedom of lepton, Higgs doublets and Higgs triplet and right-handed neutrinos, respectively, and $z = \frac{M_{\Delta_1}}{T}$, with M_{Δ_1} being the mass of decaying particle. The co-moving entropy density is given by $s = (\frac{2\pi^2}{45}) g_\star T^3$, $\zeta(3) \approx 1.202$, and $K_i(z)$ are the modified Bessel functions of type i. Let us define the total co-moving number density of the triplet be given by $\Sigma_{\Delta i} = \frac{n_{\Delta_i} + n_{\bar{\Delta}_i}}{s}$. The asymmetric densities of the particles are given by $Y_x = \frac{n_x - n_{\bar{x}}}{s}$. Where, n_x are the number density of the corresponding particles. Hence we can write the Boltzmann equations for different reaction densities to show the evolution of asymmetries Y_{Δ_1} (triplet), Y_h (Higgs), Y_l (leptons), and are given in [38, 62]:

$$szH(z)\frac{\mathrm{d}\Sigma_{\Delta_1}}{\mathrm{d}z} = -\left(\frac{\Sigma_{\Delta_1}}{\Sigma_{\Delta_1}^{eq}} - 1\right)\gamma_D - 2\left(\frac{\Sigma_{\Delta_1}^2}{\Sigma_{\Delta_1}^{eq\,2}} - 1\right)\gamma_A, \tag{52}$$

$$szH(z)\frac{\mathrm{d}Y_{\Delta_1}}{\mathrm{d}z} = -\gamma_D\left(\frac{Y_{\Delta_1}}{\Sigma_{\Delta_1}^{eq}} + \sum B_l \frac{Y_l}{Y_l^{eq}} - B_h \frac{Y_h}{Y_h^{eq}}\right), \tag{53}$$

$$szH(z)\frac{\mathrm{d}Y_h}{\mathrm{d}z} = 2\gamma_D\left(\sum B_l \epsilon_l - B_h \epsilon_h\right) - 2B_h\gamma_D\left(\frac{Y_h}{Y_h^{eq}} - \frac{Y_{\Delta_1}}{\Sigma_{\Delta_1}^{eq}}\right) - 2\left(\frac{Y_h}{Y_h^{eq}} + 2\frac{Y_l}{Y_l^{eq}}\right)(2\gamma_{ll}^{hh} + \gamma_{lh}^{\bar{l}\bar{h}}), \tag{54}$$

$$szH(z)\frac{\mathrm{d}Y_l}{\mathrm{d}z} = 2\gamma_D\left(\sum B_l\epsilon_l - B_h\epsilon_h\right) - 2B_l\gamma_D\left(\frac{Y_l}{Y_l^{eq}} + \frac{Y_{\Delta_1}}{\Sigma_{\Delta_1}^{eq}}\right) - 2\left(\frac{Y_h}{Y_h^{eq}} + 2\frac{Y_l}{Y_l^{eq}}\right)(2\gamma_{ll}^{hh} + \gamma_{lh}^{\bar{l}\bar{h}}). \tag{55}$$

The decays and inverse decays are important in the Boltzmann equations to contribute to the lepton asymmetry, which are given by

$$\gamma_D = s\Gamma_{\Delta_1}\Sigma_{\Delta_1}\frac{K_1(z)}{K_2(z)}, \tag{56}$$

and the s-wave contribution to the gauge scattering processes of triplets is

$$\gamma_A = \frac{M_{\Delta_1}T^3 e^{\frac{-2M_{\Delta_1}}{T}}}{64\pi^4}(9g^4 + 12g^2g_1^2 + 3g_1^4)\left(1 + \frac{3T}{4M_{\Delta_1}}\right), \tag{57}$$

where, g and g_1 are the SM gauge couplings. Along with the gauge scattering for triplets, we can also have $\Delta L = 2$, lepton number violating interactions. But we can safely neglect them in this case, as those will be suppressed by the heavy mass of mediating particle (Δ_1). These lepton number violating scattering processes are given as following

$$\hat{\sigma}_{ll}^{hh} = 64B_hB_l\left(\frac{\Gamma_\Delta}{M_{\Delta_1}}\right)^2\frac{s_{\mathrm{cm}}/M_{\Delta_1}^2}{(s_{\mathrm{cm}}/M_{\Delta_1}^2 - 1)^2 + \left(\frac{\Gamma_\Delta}{M_{\Delta_1}}\right)^2}, \tag{58}$$

$$\hat{\sigma}_{lh}^{\bar{l}\bar{h}} = 64B_hB_l\left(\frac{\Gamma_\Delta}{M_{\Delta_1}}\right)^2\frac{M_{\Delta_1}^2}{s_{\mathrm{cm}}}\left(\ln\left(1 + \frac{s_{\mathrm{cm}}}{M_{\Delta_1}^2}\right) - \frac{s_{\mathrm{cm}}/M_{\Delta_1}^2}{1 + s_{\mathrm{cm}}/M_{\Delta_1}^2}\right), \tag{59}$$

$$\gamma_s = \frac{M_{\Delta_1}^4}{64\pi^4}\int_{s_{\mathrm{cm\,min}}}^{\infty} ds_{\mathrm{cm}}\sqrt{s_{\mathrm{cm}}}\frac{K_1(\sqrt{s_{\mathrm{cm}}}z)\hat{\sigma}_s}{z}. \tag{60}$$

Hence from figure 9 one can see that a lepton asymmetry of order $\approx 10^{-9}$ can be achieved by using the Boltzmann equation in high mass regime of the scalar triplet. This can be converted to the baryon asymmetry during sphaleron transition with a fraction of $Y_B = aY_L$, where, $a = -\frac{8N_F + 4N_H}{22N_F + 13N_H} \approx -0.34$. N_H and N_F are the number of Higgs and fermion generations, respectively.

4.2. Low scale leptogenesis: $M_\Delta \approx \mathcal{O}(2)$ TeV

High scale leptogenesis with such a heavy triplet is very difficult to have any experimental signature in the near future. But efforts have been made to bring down the scale of leptogenesis with resonance effect, which can also explain the current neutrino oscillation data [63–68]. This low scale is not only phenomenologically viable but also can be verified by the collider experiments. We discuss the impact of leptogenesis in TeV scale along-with the constraints on coupling from the neutrino masses. We also put light on the contribution of TeV scale triplets and right-handed neutrinos to the muon $g - 2$ anomaly and lepton flavor violating (LFV) rare decays.

Constraint on couplings from out of equilibrium decay and neutrino mass. To generate nonzero asymmetry as per the Sakharov's conditions the decay rate should be less than the Hubble expansion

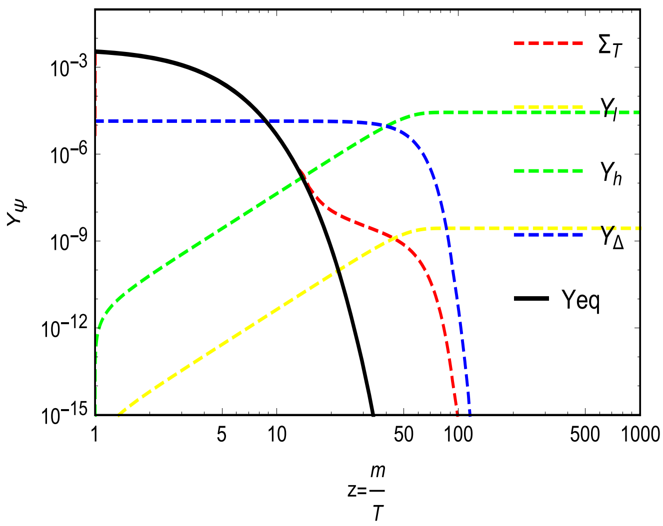

Figure 9. Figure shows the abundances of particles when lepton asymmetry is generated solely from scalar triplet. Even though the lepton asymmetry from the lepton loop vanishes, still a large asymmetry can be generated from scalar and right-handed neutrino loop.

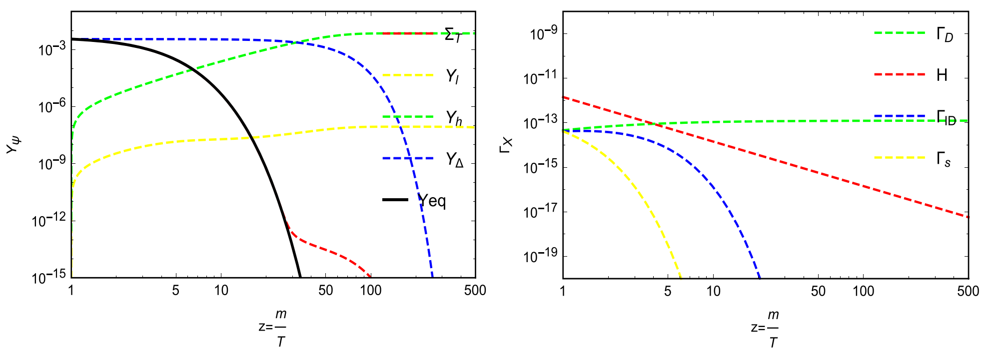

Figure 10. Left panel shows the abundances of particles when lepton asymmetry is generated solely from the TeV scale scalar triplet. Right panel represents the variation of the decay of heavy triplet and lepton number violating scattering with z.

$$\Gamma_{\Delta_1} < H \approx 1.66 \times \sqrt{g_\star}\frac{T^2}{1.2 \times 10^{19}} \text{ GeV}$$
$$< 1.38 \times 10^{-18}[T^2]_{T=M_{\Delta_1}}(g_\star \approx 100) \text{ GeV}. \tag{61}$$

From type I:

$$\sum m_\nu^{\mathrm{I}} = \frac{4y_{\nu_1}^2 v^2 \sin^2\beta}{M_{1R}} + \frac{2y_{\nu_3}^2 v^2 \sin^2\beta}{M_{1R}} + \frac{2y_{\nu_4}^2 v^2 \cos^2\beta}{M_{3R}} \leqslant 0.05 \times 10^{-9} \text{ GeV}$$
$$\approx 0.3 y_{\nu_3}^2 \sin^2\beta + y_{\nu_4}^2 \cos^2\beta \leqslant 4.1 \times 10^{-12}. \tag{62}$$

From type II:

$$\begin{aligned}\sum_i m_\nu^{\rm II} &= \frac{(4y_{ti} + 2y_{ti}{}')\mu_{iL} v^2}{M_{\Delta i}{}^2} \leqslant 0.05 \times 10^{-9}\ \text{GeV}\\ &\approx (2y_{ti} + y_{ti}{}')\mu_{iL} \leqslant 8 \times 10^{-10}\ \text{GeV}\ \ (\text{for}\ \ i = 1),\\ &\leqslant 8.3 \times 10^{-8}\ \text{GeV}\ \ (\text{for}\ \ i = 2).\end{aligned} \tag{63}$$

As the CP asymmetry turns out to be very small, another way to realize the leptogenesis is through resonance enhancement. If both of the triplets will be quasi degenerate in mass, there will be large enhancement in CP asymmetry. The expression of CP asymmetry can be written as [39]

$$\epsilon_{CP} = \frac{1}{2\pi}\frac{\text{Im}\,[\mathcal{Y}_{\Delta_1}\mathcal{Y}^\dagger_{\Delta_2}\mu^\star_{\Delta_1}\mu_{\Delta_2}]}{M^2_{\Delta_1}(\mathcal{Y}^\dagger_{\Delta_1}\mathcal{Y}_{\Delta_1})_{ii} + \mu^2_{\Delta_1}}\frac{M^2_{\Delta_1}\,\delta M^2_{12}}{(\delta M^2_{12})^2 + M^2_{\Delta_1}\Gamma^2_{\Delta_2}}, \tag{64}$$

where, $\delta M^2_{12} = M^2_{\Delta_2} - M^2_{\Delta_1}$. With the standard assumption of resonance case, provided in the literature [63], $\delta M^2_{12} \approx M_{\Delta_1}\Gamma_{\Delta_2} \ll M^2_{\Delta_1}$, we can have the self energy contribution

$$\frac{M^2_{\Delta_1}\,\delta M^2_{12}}{(\delta M^2_{12})^2 + M^2_{\Delta_1}\Gamma^2_{\Delta_2}} \approx \frac{M_{\Delta_1}}{\Gamma_{\Delta_2}} \gg 1. \tag{65}$$

Hence there will be resonance enhancement in the self energy, which contributes maximally to the CP asymmetry. With such a large CP asymmetry ($\approx\mathcal{O}(1)$), the observed baryon asymmetry can be explained by solving the Boltzmann equations as in the previous case. But at low energy the washout processes will dominantly contribute and the desired lepton asymmetry can be achieved. The Boltzmann equations for the TeV scale are solved to yield a lepton asymmetry of required order, as shown in figure 10. Some sample benchmark points for the coupling simultaneously satisfying the neutrino mass and leptogenesis are provided in table 4.

4.2.1. Comments on LFV decays and muon g-2 anomalies. LFV decay processes have received great attention in recent times which are very rare to be observed experimentally [69–74]. Efforts are being made by many experiments to look in this direction and some of them have provided a stringent upper limits on these decays. In this context, $\mu \to e\gamma$ looks to be an important process to be measured with less background from the observation point of view. The current experimental limit on this decay is Br $(\mu \to e\gamma) < 4.2 \times 10^{-13}$ from MEG collaboration [75]. In the framework of low scale leptogenesis, we can have extra contribution to rare decays $l_\alpha \to l_\beta\gamma$ due to the presence of right-handed neutrinos and Higgs. Due to the diagonal structure of scalar triplet-Yukawa coupling, there will not be any contribution to LFV from the triplet sector but we can still have it from right-handed neutrino and heavy Higgs loop. The branching ratio for this decay is given by [76]

$$Br(l_\alpha \to l_\beta\gamma) = \frac{3(4\pi)^3\alpha}{4G_F{}^2}|A_D|^2 \times Br(l_\alpha \to l_\beta\nu_\alpha\bar{\nu_\beta}). \tag{66}$$

where, $G_F \approx 10^{-5}\,\text{GeV}^{-2}$ is the Fermi constant and α is the electromagnetic fine structure constant and A_D is the dipole contribution, which is given by

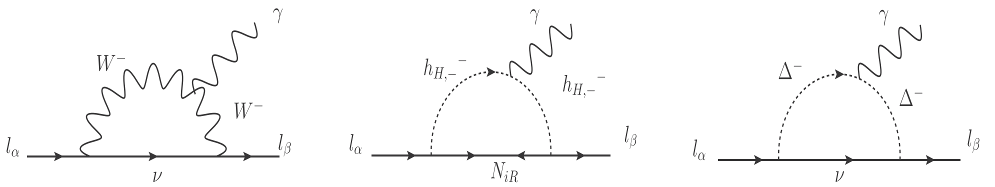

Figure 11. Feynman diagrams represent the Lepton flavor violating rare decays and muon g-2 anomaly in one loop level.

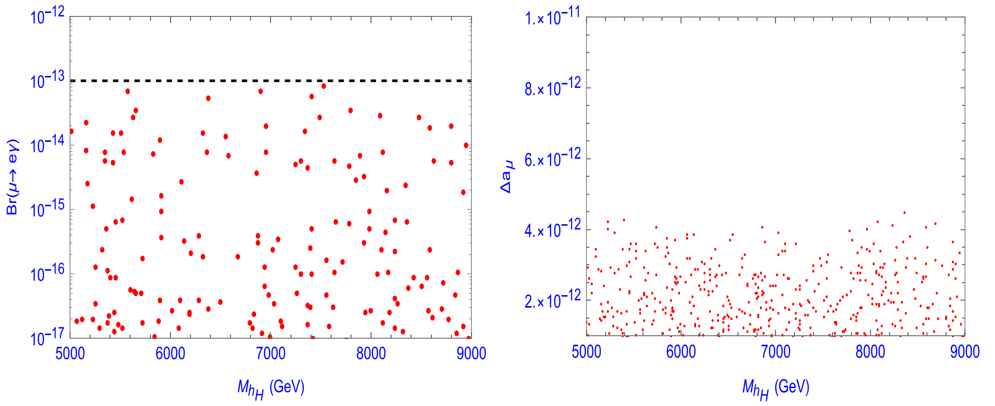

Figure 12. The left most panel shows the parameter space of heavy Higgs mass allowed by the branching of lepton flavor violating decay $\mu \to e\gamma$, which is coming in order of less than 10^{-13} as per the experimental bound. The right panel shows the variation of same Higgs mass that gives the viable range of muon anomalous magnetic moment allowed by experiment.

$$A_D = \sum_i \frac{(\mathcal{Y}^{\nu}_{H,-})_{\alpha i}(\mathcal{Y}^{\nu\star}_{H,-})_{\beta i} f(x)}{2(4\pi)^2 M^2_{h_H}}. \tag{67}$$

Here $\mathcal{Y}^{\nu}_{H,-}$'s are the Yukawa coupling matrices in equation (40), M_{h_H} is the mass of heavy Higgs (H_H, H_-) and $f(x)$ is the loop function, with $x = \frac{M^2_{iR}}{M^2_{h_H}}$, and is given by

$$f(x) = \frac{1 - 6x + 3x^2 + 2x^3 - 6x^2 \log x}{6(1-x)^4}. \tag{68}$$

When we have $\alpha = \beta$, the above diagrams will contribute to the muon anomalous magnetic moment as

$$\delta a_\mu = \frac{1}{16\pi^2}\left[\frac{m^2_\mu}{M^2_{h_H}}\sum_i |(\mathcal{Y}^{\nu}_{H,-})_{\mu\mu}|^2 f(x) + \frac{m^2_\mu}{M^2_{\Delta_i}}\sum_i |y_{ti}|^2 f(x_t)\right], \tag{69}$$

where, $f(x_t)$ will follow the same expression of $f(x)$, with $x_t = \frac{m^2_\nu}{M^2_{\Delta_i}}$.

From the left panel of figure 12, one can see that there will be extra contribution to rare LFV decay $\mu \to e\gamma$ from the right-handed neutrinos and heavy Higgs mediated diagram (shown in the middle panel of figure 11) with a mass scale of order $\mathcal{O}(10)$ TeV. The right panel of figure 12 represents the appreciable contribution to muon g-2 anomaly due to the presence of extra particles in the model shown in the right panels of figure 11.

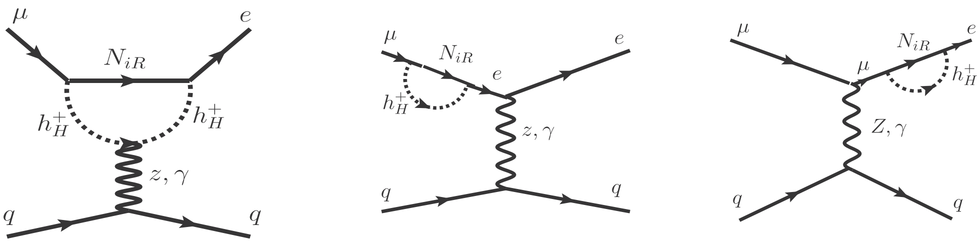

Figure 13. Feynman diagrams contribute to μ to e conversion in nucleus.

4.2.2. Comment on μ to e conversion in the nucleus. The most stringent constraint on LFV decays preferred by the $\mu \to e\gamma$, however, the improved sensitivity is expected from the μ to e conversion in the nucleus in coming decades. Various projects like Mu2e, DeeMe, COMET and PRISM/PRIME [77–79] are on its peak to reach a bound from current limit 4.3×10^{-14} (for Ti Nucleus) to future sensitivity up to 10^{-18}. For models with dominant contribution from dipole operators, the μ to e conversion is little suppressed as compared to the $\mu \to e\gamma$. But this is not a rigid case and we will compare the contribution from both LFV decays and μ to e conversion in the nucleus (figure 13) in this model

$$\begin{aligned} CR(\mu \to e,\, nucleus) = & \frac{P_e E_e m_\mu^3 G_f^2 \alpha_{\rm em}^3 Z_{eff}^4 F_p^2}{8\pi^2 Z \Gamma_{\rm capt}} \\ & \times [|(Z+N)(g_{LV}^0 + g_{LS}^0) + (Z-N)(g_{LV}^1 + g_{LS}^1)|^2 \\ & + |(Z+N)(g_{RV}^0 + g_{RS}^0) + (Z-N)(g_{RV}^1 + g_{RS}^1)|^2]. \end{aligned} \tag{70}$$

Here, Z and N are the atomic number and number of neutrons, $Z_{\rm eff}$ is the effective charge of the nucleus and $\Gamma_{\rm capt}$ is the total muon capture rate. The definition and values of all the parameters in the above expression are given in [80–89]. The loop contribution for the non-dipole component of photonic loop is given by

$$G(x) = \frac{2 - 9x + 18x^2 - 11x^3 + 6x^3 \log x}{6(1-x)^4}, \quad x = \frac{M_{iR}^2}{M_{h_H}^2}. \tag{71}$$

From figure 14, we found the μ to e conversion in the Ti nucleus is compatible with the parameter space of the model and contribution is dominated by the dipole and non dipole component of the photonic loop. The non dipole contribution is a little larger than the dipole contribution, when right-handed neutrino mass is much larger than the heavy Higgs mass.

Finally commenting on the collider searches of the triplet scalar, where, the processes important for the production of doubly or singly charged Higgs are as following

$$\begin{aligned} & p, p \to Z/\gamma \to \Delta_i^{++}\Delta_i^{--} \\ & p, p \to Z/\gamma \to \Delta_i^{+}\Delta_i^{-} \\ & p, p \to W^+ \to \Delta_i^{++}\Delta_i^{-}. \end{aligned} \tag{72}$$

As the pair produced doubly charged scalar which decay to two equal signed leptons, hence can be obtained in the collider with four lepton ($l^+l^+l^-l^-$) signature in final states (figure 15). Various studies in the literature are focused on the constraint on scalar mass depending on the final state leptons [90–98]. For an example, the mass of the doubly charged

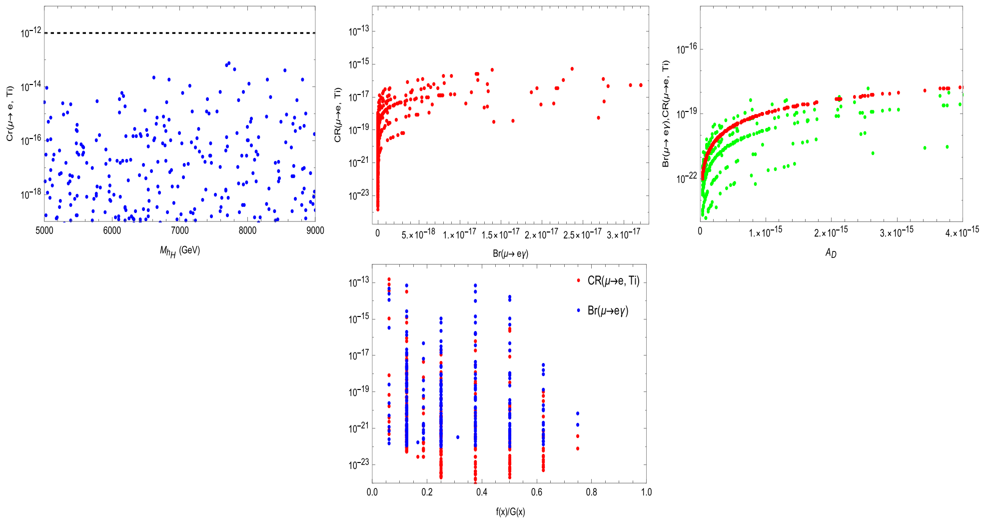

Figure 14. The left most panel shows the variation of heavy Higgs mass with the conversion ratio of $\mu \rightarrow e$ in Titanium nucleus, which is coming in order of less than 10^{-12} as per the experimental bound. The middle panel shows the variation of $\mu \rightarrow e\gamma$ branching with the μ to e conversion. Extreme right panel represents the variation of dipole operator with $\mu \rightarrow e\gamma$ branching (red) and μ to e conversion ratio (green). The lower panel represents the variation of loop functions for dipole (f) and non-dipole (G) contribution with the μ to e branching and conversion factor.

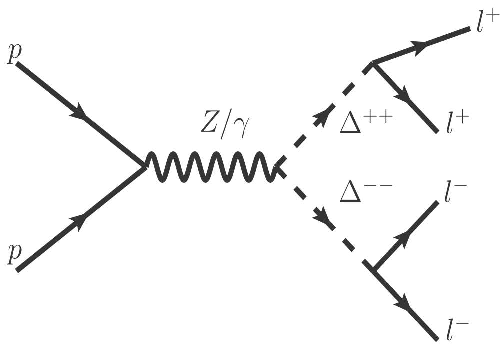

Figure 15. Production and decay of doubly charged scalar to four lepton final states.

scalar is found to be >943 GeV with the current LHC limit with 36 fb^{-1} luminosity and 14 TeV energy. The HL-LHC after up gradation with an integrated luminosity of 3ab^{-1} and same energy is predicted to search a doubly charged scalar mass greater than 2.5 TeV, whereas the highly improved HE-LHC with 27 TeV available energy can search for a Higgs mass larger than 4.9 TeV [99]. Hence considering the limits, in our model the mass of the triplet is greater than 1.6 TeV from the leptogenesis constraint and hence is more sensitive to HL-LHC. We used CalcHEP to get the production cross section of doubly charged Higgs, which is found to be order of 10^{-5}pb for charged scalar mass of order 1 TeV, as shown in figure 16. The detailed study of collider is beyond the scope of this work.

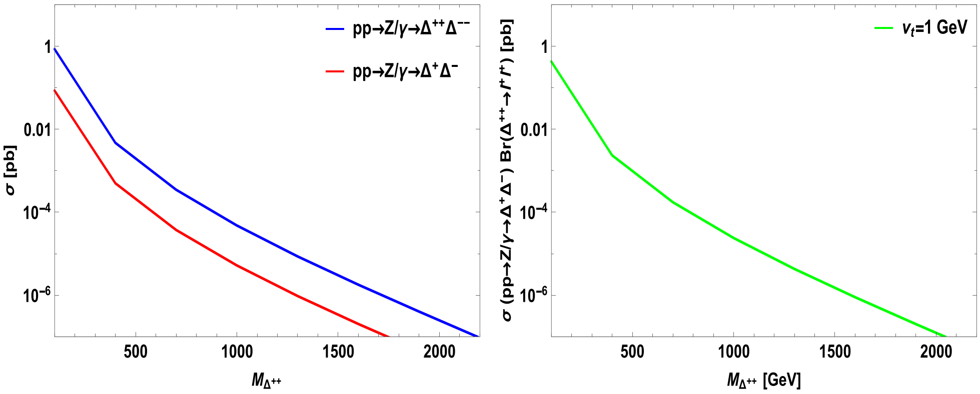

Figure 16. Production of doubly and singly charged Higgs in the pp-collision at LHC is shown in the left panel and the decay of charged scalars to four lepton final state is represented in the right panel with consideration of 50% branching ratio.

5. Summary

In this article, an attempt has been made to understand the lepton asymmetry using the simplest discrete symmetry S_3 along with the Z_2 symmetry. Here, we first considered the type-II seesaw mechanism by introducing the scalar triplets but the triplet-lepton Yukawa matrix turned out to be diagonal. So in order to explain the neutrino masses and mixing, we included additional right-handed neutrinos in the type-I seesaw scenario. We considered the combination of both type-I and type-II seesaw scenarios to accommodate current results in the neutrino sector, so also attempted to explain the observed baryon asymmetry. Due to the presence of both right-handed neutrinos and scalar triplets we considered two different scenarios, in which the first case we discussed the scenario where the scalar triplets are lighter than the right-handed neutrino masses and explained the lepton asymmetry in high mass regime. Thereafter, we considered the case where the masses of RHNs and scalar triplets are much smaller (of the order of TeV), which can be tested in future colliders [90–99], and explained the leptogenesis by resonance enhancement from self energy loop with quasi degenerate triplets. We used the Boltzmann equation for the first case and discussed the results, which showed that the combined effect of type-I+II enhances the CP asymmetry. There is no contribution from the triplet mediated process but due to the presence of right-handed neutrinos and heavy Higgs there will be contributions to the $\mu \to e\gamma$ process and hence one can explain the available LFV result. Moreover, in this model for the muon anomalous magnetic moment, there will be contributions from the type I and II sectors. Hence the current framework is more interesting as it not only explains the neutrino mass and leptogenesis results but also satisfies the experimental bounds on LFV branching ratio and muon anomalous magnetic moment simultaneously in the low mass regime with quasi degenerate triplets. Interestingly, the low mass regime also opens up the exciting possibilities of scalar triplets to be tested in future collider experiments.

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Appendix

S_3 group includes three-dimensional reducible representation which can be reduced to a doublet and a singlet (i.e. $3 = 2 \oplus 1$). If $(a_1 \quad a_2)$ and $(b_1 \quad b_2)$ transform as doublets under S_3, the tensor product rules are summarized as [12]

$$(a_1 \;\; a_2)_2^T \otimes (b_1 \;\; b_2)_2^T = \begin{pmatrix} a_1b_2 + a_2b_1 \\ a_1b_1 - a_2b_2 \end{pmatrix}_2 \oplus (a_1b_1 + a_2b_2)_1 \oplus (a_1b_2 - a_2b_1)_{1'}.$$

$$(a_1 \;\; a_2)_2^T \otimes (b')_1 = (a_1b' \;\; a_2b')_2^T, \qquad (a_1 \;\; a_2)_2^T \otimes (b')_{1'} = (-a_2b' \;\; a_1b')_2^T.$$

$$(a)_1 \otimes (b)_{1'} = (ab)_{1'}, \qquad (a)_{1'} \otimes (b)_{1'} = (ab)_1.$$

ORCID iDs

Subhasmita Mishra https://orcid.org/0000-0002-0593-9127
Anjan Giri https://orcid.org/0000-0002-8895-0128

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