# $\mu-\tau$ Reflection Symmetry in in Lepton Mixing and Radiatively Generated Leptogenesis

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We study a $\mu-\tau$ reflection symmetry in neutrino sector realized at the GUT scale in the context of the seesaw model. In our scenario, the exact $\mu-\tau$ reflection symmetry realized in the basis where the charged lepton and heavy Majorana mass matrices are diagonal, leads to vanishing lepton asymmetries. We find that, in the minimal supersymmetry extension of the seesaw model with appropriate values of $\tan\beta$, the renormalization group (RG) evolution from the GUT scale to seesaw scale can induce a successful leptogenesis. It is shown that the right amount of the baryon asymmetries $\eta_B$ can be achieved via so-called resonant leptogenesis, which can be realized at rather low seesaw scale in our scenario, so that the well-known gravitino problem is safely avoided. In this work, we consider both flavor dependent and flavor independent leptogenesis, and demonstrate how they lead to different amounts of baryon asymmetries in detail.

## 1. Introduction

Recently, several neutrino mass matrix textures have been proposed that predict a maximal value for the Dirac CP violating phase $\delta_{CP}=\pi/2$ and maximal atmospheric mixing $\theta_{23}=\pi/4$. Our scenario is constructed in the context of a seesaw model, and the right-handed neutrinos transform into the charge conjugates of themselves, $N_i \rightarrow N_i^c$, and the neutrino Dirac Yukawa couplings is given as Eq. (3) under the $\mu-\tau$ reflection symmetry (MTRS) which is suggested in Ref. [1]. After seesawing, the effective mass matrix represents the $\mu-\tau$ reflection symmetry which means $\delta_{CP}=\pi/2$ and $\theta_{23}=\pi/4$ and predicts the Jarlskog invariant to be maximal.

The symmetry is strongly broken in the charged lepton sector. As a result, we expect the radiative corrections induced by the charged lepton sector to break this symmetry in the neutrino sector, too. In this work we would like to propose that the precise MTRS, imposed in Ref. [1], exists only at high energy scale such as GUT scale and a renormalization group (RG) evolution from high scale to low scale gives rise to the breaking of MTRS in the lepton sector without introducing any ad hoc soft symmetry breaking terms. In this work we shall show that such RG effecs in supersymmetric seesaw model can lead to successful leptogenesis which is absent in the exact MTRS model. We consider both flavor-dependent and flavor-independent leptogenesis and show how they lead to different amounts of lepton asymmetries in detail.

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## 2. Suppersymmetric Seesaw Model with MTRS realized at GUT scale

Let us begin by considering a supersymmetric version of the seesaw model, which is given as the following leptonic superpotential:

$$W_l = \widehat{l}_L^c \mathbf{Y}_l \widehat{L}\cdot\widehat{H}_1 + \widehat{N}_L^c \mathbf{Y}_\nu \widehat{L}\cdot\widehat{H}_2 - \frac{1}{2}\widehat{N}_L^{cT}\mathbf{M}_R\widehat{N}_L^c \ . \quad (1)$$

After spontaneous symmetry breaking, the seesaw mechanism leads to a following effective light neutrino mass term,

$$m_{\rm eff} = -\mathbf{Y}_\nu^T \mathbf{M}_R^{-1}\mathbf{Y}_\nu \upsilon_2^2. \quad (2)$$

Suppose that under the MTRS, the right-handed neutrinos transform into the charge conjugates of themselves, $N_i \rightarrow N_i^c$. In order for the right-handed neutrino mass matrix to be symmetric under the $\mu-\tau$ reflection, it should be real and then the neutrino Dirac Yukawa matrix and the heavy Majorana neutrino mass matrix are given as

$$\begin{aligned} \mathbf{Y}_\nu &= b_3 \begin{pmatrix} \lambda & \omega e^{i\phi_1} & \omega e^{-i\phi_1} \\ \chi & \kappa e^{i\phi_2} & \kappa e^{-i\phi_2} \\ \rho & e^{i\phi_3} & e^{-i\phi_3} \end{pmatrix}, \\ \mathbf{M}_R &= \begin{pmatrix} M_1 & 0 & 0 \\ 0 & M_2 & 0 \\ 0 & 0 & M_3 \end{pmatrix}, \end{aligned} \quad (3)$$

and all the lepton flavor mixing is in $\mathbf{Y}_\nu$, where the parameters of the neutrino Dirac Yukawa matrix are all real. We assume an exact degenerate of two heavy Majorana neutrinos so that $M_1 = M_2 \equiv M \ll M_3$ and that the right-handed neutrino $N_3$ dominates in generation of the light mass matrix (single right-handed neutrino dominant) [2, 3], then we can left with $\rho = 0$ at GUT scale.

The CP-violation in neutrino oscillations is determined proportional to

$$J_{CP} = \frac{1}{8}\sin 2\theta_{12}\sin 2\theta_{23}\sin 2\theta_{13}\cos\theta_{13}\sin\delta_{CP} \quad (4)$$

due to the MTRS, the CP-violation is maximal. On the other hand, a non-zero leptonic asymmetry can be generated if and if only the $CP$-odd invariants $\tilde{J}_{CP} = \mathrm{ImTr}[HM_R^\dagger M_R M_R^\dagger H^T M_R]$ does not vanish [4, 5]. Since $\tilde{J}_{CP}$ can be expressed in the form

$$\begin{aligned} \tilde{J}_{CP} &= \sum_{i<j}\{M_i M_j (M_j^2 - M_i^2)\mathrm{Im}[H_{ij}^2]\} \neq 0, \\ H &= \mathbf{Y}_\nu \mathbf{Y}_\nu^\dagger. \end{aligned} \quad (5)$$

In order to accommodate leptogenesis, it requires not only $M_i \neq M_j$ but also $\mathrm{Im}[H_{ij}^2] \neq 0, (i \neq j = 1, 2, 3)$, at the leptogenesis scale. We can see that, the Hermitian parameter $H$ in the limit of exact $MTRS$ leads to vanishing lepton asymmetry which is disastrous for successful leptogenesis, see

$$H \equiv b_3^2 \begin{pmatrix} \lambda^2 + 2\omega^2 & \lambda\chi + 2\omega\kappa c_{12} & 2\omega c_{13} \\ \lambda\chi + 2\omega\kappa c_{12} & \chi^2 + 2\kappa^2 & 2\kappa c_{23} \\ 2\omega c_{13} & 2\kappa c_{23} & 2 \end{pmatrix}, \quad (6)$$

where $c_{ij} \equiv \cos\Delta\phi_{ij}$. To generate non-vanishing lepton asymmetry, we need to break the exact MTRS texture of $Y_\nu$ proposed above. In our scenario, as will be shown later, only the RG evolution is responsible for such a breaking required for leptogenesis.

## 3. Confronting with low-energy neutrino data

Before discussing how to achieve leptogenesis in MTRS model, we first examine if it is consistent with low-energy neutrino data. To do so, we need renormalization group evolution of neutrino Dirac-Yukawa matrix $Y_\nu$ and heavy neutrino mass matrix $M_R$ with MTRS from the GUT scale to the seesaw scale by numerically solving the relevant RGE's represented in Ref. [6].

In our numerical calculation of the RG running effects, we first fix the values of heavy Majorana neutrino masses with hierarchy $M_3 \gg M_1 = M_2 = M$ and $\tan\beta$, then we solve the RGE's by varying input values of all the parameter space $\{b_3, \chi, \kappa, \omega, \lambda, \phi_1, \phi_2\}$ given at GUT scale (here we put $\phi_3 = 0$ since, as will be shown later, lepton asymmetries do not depend on $\phi_3$). Then finally we determine the parameter space allowed by low energy neutrino experimental data. In our numerical calculation, we have five experimental results for neutrino mixing parameters and mass squared differences given at $3\sigma$ by [7],

$$\begin{aligned} &29.3^\circ < \theta_{12} < 39.2^\circ, \quad 35.7^\circ < \theta_{23} < 55.6^\circ, \\ &0^\circ < \theta_{13} < 11.5^\circ, \quad 7.1 < \Delta m_{21}^2[10^{-5}\ \mathrm{eV}^2] < 8.9, \\ &2.0 < \Delta m_{31}^2[10^{-3}\ \mathrm{eV}^2] < 3.2. \end{aligned} \quad (7)$$

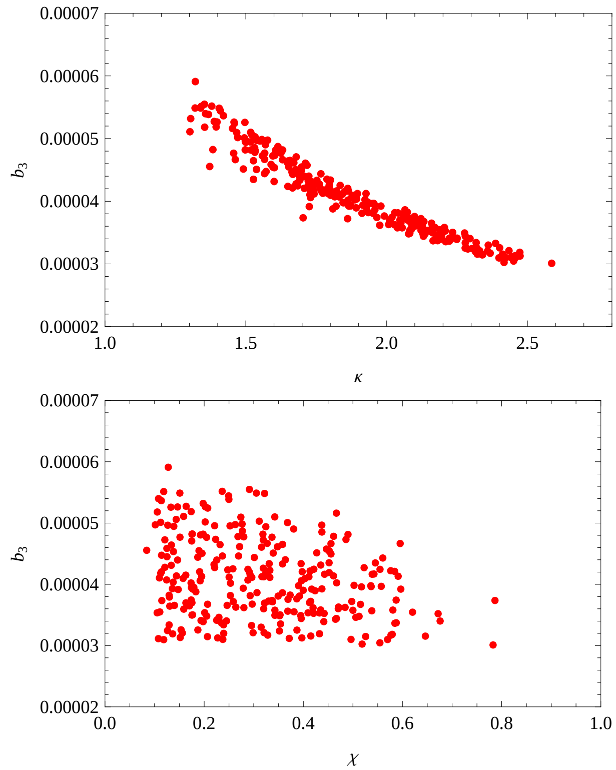


Figure 1: The allowed regions of the parameters $b_3, \kappa, \chi$.

In figures 1 and 2, we shown the parameter regions constrained by the experimental data given in Eq. (7). The figures exhibit how the parameters $\{b_3, \kappa, \chi, \lambda, \omega\}$ are correlated. Here we adopted $M = 5 \times 10^6$Gev, $M_3 = 10^{12}$GeV and $\tan\beta = 10$ as inputs.

Figure 3 shows the correlation between the parameter $b_3$ and the phase difference $\Delta\phi_{12} = \phi_2 - \phi_1$, among the high-energy CP phases in the neutrino Dirac-Yukawa coupling matrix $Y_\nu$ given in Eq. (3). We find that the allowed regions of $\Delta\phi_{12}$ are $40^0 \leq \Delta\phi_{12} \leq 140^0$ and $220^0 \leq \Delta\phi_{12} \leq 320^0$.

## 4. Relevant RGE's in MSSM

In the minimal supersymmerty standard model (MSSM), the radiative behavior of the Dirac neutrino Yukawa matrix can written as [8]

$$\frac{dY_\nu}{dt} = Y_\nu[(T - 3g_2^2 - \frac{3}{5}g_1^2) + (Y_l^\dagger Y_l + 3Y_\nu^\dagger Y_\nu)], \quad (8)$$

where $t = \frac{1}{16\pi^2}\ln(Q/Q_{GUT})$ with an arbitrary renormalization scale $Q$; $T = Tr(3Y_u^\dagger Y_u + Y_\nu^\dagger Y_\nu)$, $Y_u$ and $Y_l$ are the Yukawa matrices for up-type quarks and

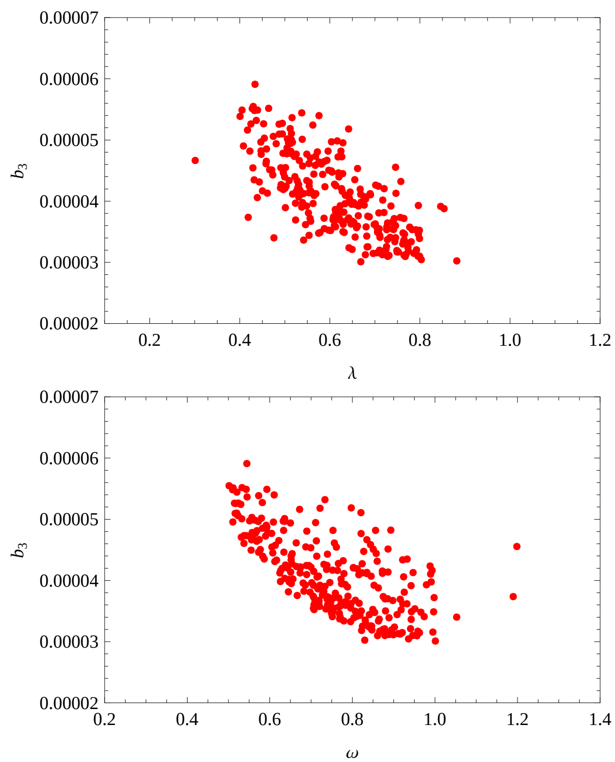


Figure 2: The allowed regions of the parameters $\lambda, \omega$.

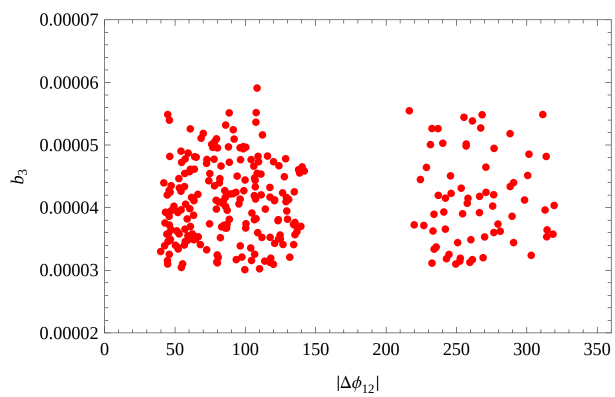


Figure 3: The correlation between the parameter $b_3$ and and the phase difference $\Delta\phi_{12}$, among the high-energy CP phases in the Dirac-Yukawa matrix.

charged leptons and $g_{2,1}$ are the $SU(2)_L$ and $U(1)_Y$ gauge coupling constants. And the RG equation for the heavy Majorana mass matrix is given by

$$\frac{dM_R}{dt} = 2[(Y_\nu Y_\nu^\dagger)M_R + M_R(Y_\nu Y_\nu^\dagger)^T]. \qquad (9)$$

For latter convenience, let us reformulate the RG equation (9) in the basis where $M_R$ is diagonal. Since $M_R$ is symmetric, it can be diagonalized by a unitary matrix $V$,

$$V^T M_R V = \mathrm{Diag.}(M_1, M_2, M_3). \qquad (10)$$

Note that, as the structure of the mass matrix $M_R$ change with the evolution of the scale, the unitary matrix $V$ depends on the scale, too. The RG evolution of the matrix $V(t)$ can be written as

$$\frac{d}{dt}V = VA, \qquad (11)$$

where $A$ is an anti-Hermitian matrix $A^\dagger = -A$ due to the unitary of $V$. Then, differentiating Eq. (10), we obtain

$$\begin{aligned}\frac{dM_i\delta_{ij}}{dt} &= A_{ij}^T M_j + M_i A_{ij} + 2\{V^T[(Y_\nu Y_\nu^\dagger)M_R \\ &\quad + M_R(Y_\nu Y_\nu^\dagger)^T]V\}_{ij}. \qquad (12)\end{aligned}$$

Absorbing the unitary transformation into the Dirac-Yukawa coupling $Y_\nu \equiv V^T Y_\nu$, the real diagonal part of Eq. (12) becomes

$$\frac{dM_i}{dt} = 4M_i(Y_\nu Y_\nu^\dagger)_{ii}. \qquad (13)$$

On the other hand, the off diagonal part of Eq. (12) leads to

$$\begin{aligned}A_{ij} &= 2\frac{M_j + M_i}{M_j - M_i}\mathrm{Re}[(Y_\nu Y_\nu^\dagger)_{ij}] \\ &\quad + 2i\frac{M_j - M_i}{M_j + M_i}\mathrm{Im}[(Y_\nu Y_\nu^\dagger)_{ij}], \quad (i \neq j). \qquad (14)\end{aligned}$$

Note that the real part of $A_{ij}$ is singular for the degenerate cases with $M_i = M_j$. The singularity in $\mathrm{Re}[A_{ij}]$ can be eliminated by an appropriate rotation between degenerate heavy Majorana neutrino states. Such a rotation does not change any physics and it is equivalent to absorbing the rotation matrix into Dirac-Yukawa matrix $Y_\nu$, $Y_\nu \longrightarrow \widetilde{Y}_\nu = RY_\nu$, where the matrix $R$, in our case, particular rotating 1st and 2rd generations of heavy Majorana neutrinos chich can be parameterized as

$$R(\alpha) = \begin{pmatrix} \cos\alpha & \sin\alpha & 0 \\ -\sin\alpha & \cos\alpha & 0 \\ 0 & 0 & 1 \end{pmatrix}. \qquad (15)$$

The rotating angle $\alpha$ is obtained from the condition $\mathrm{Re}[\widetilde{H}_{ij}] = 0$, where $\widetilde{H} = \widetilde{Y}_\nu \widetilde{Y}_\nu^\dagger = RHR^T$

$$\tan 2\alpha = \frac{2H_{12}}{H_{11} - H_{22}} = \frac{2(\lambda\chi + 2\omega\kappa\cos\Delta\phi_{12})}{\lambda^2 + 2\omega^2 - \chi^2 - 2\kappa^2}. \qquad (16)$$

The running of the mass-slitting parameter $\delta_N = 1 - M_2/M_1$ which quantifies the degree of degeneracy between $M_1$ and $M_2$ can be derived from Eq. (13)

$$\frac{d\delta_N}{dt} = 4(1 - \delta_N)(\widetilde{H}_{11} - \widetilde{H}_{22}) \simeq \frac{8H_{12}}{\sin 2\alpha}, \qquad (17)$$

$$\delta_N \simeq \frac{8b_3^2}{\sin 2\alpha}(\lambda\chi + 2\omega\kappa\cos\Delta\phi_{12}).t, \qquad (18)$$

Next, we consider the running of the parameter $\widetilde{H}$ which is relevant for both flavor-dependent and flavor-independent leptogenesis in the diagonal basis of $M_R$ after rotating

$$\frac{d\widetilde{H}}{dt} = 2\Big[(T - 3g_2^2 - \frac{3}{5}g_1^2)\widetilde{H} + \widetilde{Y}_\nu(\widetilde{Y}_l^\dagger\widetilde{Y}_l)\widetilde{Y}_\nu^\dagger + 3\widetilde{H}^2\Big] + A^T H + HA^* \quad (19)$$

Then it is straightforward to obtain the radiative terms which are relevant to leptogenesis

$$\begin{aligned}\mathrm{Re}[\widetilde{H}_{12}] &= \mathrm{Re}[\widetilde{H}_{21}] \simeq y_\tau^2 b_3^2\Big[\frac{1}{2}(\kappa^2-\omega^2)\sin 2\alpha + \kappa\omega\cos 2\alpha\cos\Delta\phi_{12}\Big]\cdot t,\\ \mathrm{Im}[\widetilde{H}_{12}] &= -\mathrm{Im}[\widetilde{H}_{21}] \simeq -2y_\tau^2 b_3^2\omega\kappa\sin\Delta\phi_{12}\cdot t\end{aligned} \quad (20)$$

## 4.1. Flavor independent leptogenesis

In our scheme. we take the heavy Majorana neutrino masses hierarchy $M_3 \gg M_1 = M_2$ at GUT scale, and the degenerate of $M_1, M_2$ is broken through the RG evolution from GUT scale to seesaw scale. The lepton asymmetry required for successful leptogenesis is governed by the decay of $N_1, N_2$. And when two lightest heavy Majorana neutrinos are nearly degenerate, the CP asymmetries through their decays get dominant contributions from self-energy contribution diagrams and given by [9]

$$\begin{aligned}\varepsilon_{1(2)} &\simeq \frac{\mathrm{Im}[(Y_\nu Y_\nu^\dagger)_{12}^2]}{16\pi(Y_\nu Y_\nu^\dagger)_{11(22)}\delta_N}\Big(1+\frac{\Gamma_{2(1)}^2}{4M_{2(1)}^2\delta_N^2}\Big)^{-1}\\ &\simeq \frac{\mathrm{Im}[(Y_\nu Y_\nu^\dagger)_{12}]\cdot\mathrm{Re}[(Y_\nu Y_\nu^\dagger)_{12}]}{8\pi(Y_\nu Y_\nu^\dagger)_{11(22)}\delta_N}\\ &\simeq \frac{-\omega\kappa\sin 2\alpha\Big[\frac{1}{2}(\kappa^2-\omega^2)+\kappa\omega\cos 2\alpha\cos\Delta\phi_{12}\Big]}{32\pi\Big[\lambda\chi + 2\omega\kappa\cos\Delta\phi_{12}\Big]h_{1(2)}}\\ &\quad\times y_\tau^4\sin\Delta\phi_{12}\cdot t,\end{aligned} \quad (21)$$

where $h_{1,2}$ are defined as

$$\begin{aligned}h_1 &= \widetilde{H}_{11}/b_3^2 = (2\kappa^2+\omega^2)\sin^2\alpha + (\lambda^2+2\omega^2)\cos^2\alpha + [\lambda\chi + 2\kappa\omega\cos\Delta\phi_{12}]\sin 2\alpha,\\ h_2 &= \widetilde{H}_{22}/b_3^2 = (2\kappa^2+\omega^2)\sin^2\alpha + (\lambda^2+2\omega^2)\cos^2\alpha - [\lambda\chi + 2\kappa\omega\cos\Delta\phi_{12}]\sin 2\alpha.\end{aligned} \quad (22)$$

In the flavor independent leptogenesis, we are always in the strong washout regime, then the generated baryon asymmetry is given by [10]

$$Y_B \simeq \frac{10}{31}\sum_{i=1}^{2} 0.3\frac{\varepsilon_i}{g_*}\Big(\frac{0.55\times 10^{-3}\ \mathrm{eV}}{\widetilde{m}_i}\Big)^{1.16}, \quad (23)$$

where $g_* \equiv g_{*MSSM} \simeq 228.75$ is the effective number degrees of freedom and

$$\widetilde{m}_i \equiv \frac{\widetilde{H}_{ii}}{M_i}\upsilon_2^2. \quad (24)$$

Then, the resulting baryon-to-photon ratio becomes

$$\eta_B = [\frac{s}{n_\gamma}]_0\cdot\frac{n_B}{s} \simeq 7.0 Y_B. \quad (25)$$

## 4.2. Flavor dependent leptogenesis

Considering flavor effects, the CP asymmetries generated through the decays of $N_1, N_2$ for each flavor are given by [9]

$$\begin{aligned}\varepsilon_{1(2)}^\alpha &= \frac{\mathrm{Im}[(Y_\nu Y_\nu^\dagger)_{12}(Y_\nu)_{1(2)\alpha}(Y_\nu)^*_{2(1)\alpha}]}{16\pi(Y_\nu Y_\nu^\dagger)_{11(22)}\delta_N}\times\Big(1+\frac{\Gamma_{2(1)}^2}{4M_{2(1)}^2\delta_N^2}\Big)^{-1}\\ &\simeq \frac{\mathrm{Im}[\widetilde{H}_{12}(\widetilde{Y}_\nu)_{1(2)\alpha}(\widetilde{Y}_\nu)^*_{2(1)\alpha}]}{16\pi(\widetilde{Y}_\nu\widetilde{Y}_\nu^\dagger)_{11(22)}\delta_N}.\end{aligned} \quad (26)$$

We obtain, from the matrix $\widetilde{H}$ and Eq. (20), the relevant terms for the CP asymmetries $\varepsilon_{1(2)}^\alpha$

$$\begin{aligned}&\mathrm{Im}[(\widetilde{H})_{12}(\widetilde{Y}_\nu)_{1e}(\widetilde{Y}_\nu)^*_{2e}] = -\mathrm{Im}[(\widetilde{H})_{21}(\widetilde{Y}_\nu)_{2e}(\widetilde{Y}_\nu)^*_{1e}]\\ &\quad\simeq -2y_\tau^2 b_3^2\omega\kappa s_{12}\Big(\frac{\chi^2-\lambda^2}{2}s + \lambda\chi c\Big)\cdot t,\\ &\mathrm{Im}[(\widetilde{H})_{12}(\widetilde{Y}_\nu)_{1\mu}(\widetilde{Y}_\nu)^*_{2\mu}] = -\mathrm{Im}[(\widetilde{H})_{21}(\widetilde{Y}_\nu)_{2\mu}(\widetilde{Y}_\nu)^*_{1\mu}]\\ &\quad\simeq -y_\tau^2 b_3^2\omega\kappa s_{21}\Big(\frac{\kappa^2-\omega^2}{2}s + \kappa\omega c c_{12}\Big)\cdot t,\\ &\mathrm{Im}[(\widetilde{H})_{12}(\widetilde{Y}_\nu)_{1\tau}(\widetilde{Y}_\nu)^*_{2\tau}] = -\mathrm{Im}[(\widetilde{H})_{21}(\widetilde{Y}_\nu)_{2\tau}(\widetilde{Y}_\nu)^*_{1\tau}]\\ &\quad\simeq -3y_\tau^2 b_3^2\omega\kappa s_{12}\Big(\frac{\kappa^2-\omega^2}{2}s + \kappa\omega c c_{12}\Big)\cdot t,\end{aligned} \quad (27)$$

where $s_{12} \equiv \sin\Delta\phi_{12}, c_{12} \equiv \sin\Delta\phi_{12}, s \equiv \sin 2\alpha$ and $c \equiv \cos 2\alpha$. Then we have

$$\begin{aligned}\varepsilon_{1(2)}^e &\simeq \frac{-y_\tau^2\omega\kappa s(\frac{\chi^2-\lambda^2}{2}s + \lambda\chi c)}{64\pi(\lambda\chi + 2\omega\kappa c_{12})\cdot h_{1(2)}}s_{12},\\ \varepsilon_{1(2)}^\mu &\simeq \frac{-y_\tau^2\omega\kappa s(\frac{\kappa^2-\omega^2}{2}s + \kappa\omega c c_{12})}{128\pi(\lambda\chi + 2\omega\kappa c_{12})\cdot h_{1(2)}}s_{12},\\ \varepsilon_{1(2)}^\tau &\simeq \frac{-y_\tau^2\omega\kappa s(\frac{\kappa^2-\omega^2}{2}s + \kappa\omega c c_{12})}{128\pi(\lambda\chi + 2\omega\kappa c_{12})\cdot h_{1(2)}}s_{12}.\end{aligned} \quad (28)$$

In this work, we focus on the case where all flavor processes are in the strong washout regime, hence the lepton asymmetry in the $l_\alpha$ ($\alpha = e, \mu, \tau$) decay channel of $N_i$ decay is given [10]

$$Y_i^\alpha \simeq 0.3\frac{\varepsilon_i^\alpha}{g_*}\Big(\frac{0.55\times 10^{-3}\ \mathrm{eV}}{\widetilde{m}_i^\alpha}\Big)^{1.16} \quad (29)$$

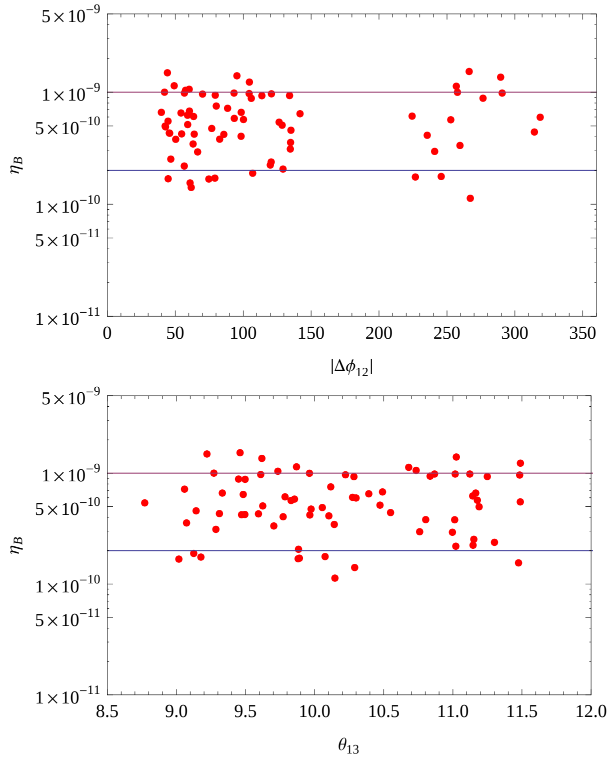


Figure 4: The prediction of flavor-dependent leptogenesis over the phase difference $\Delta\phi_{12}$ [upper panel] and the predicted of the mixing angle $\theta_{13}$ [lower panel] with $\tan\beta = 10$. The horizontal lines are the current bound of $\eta_B$.

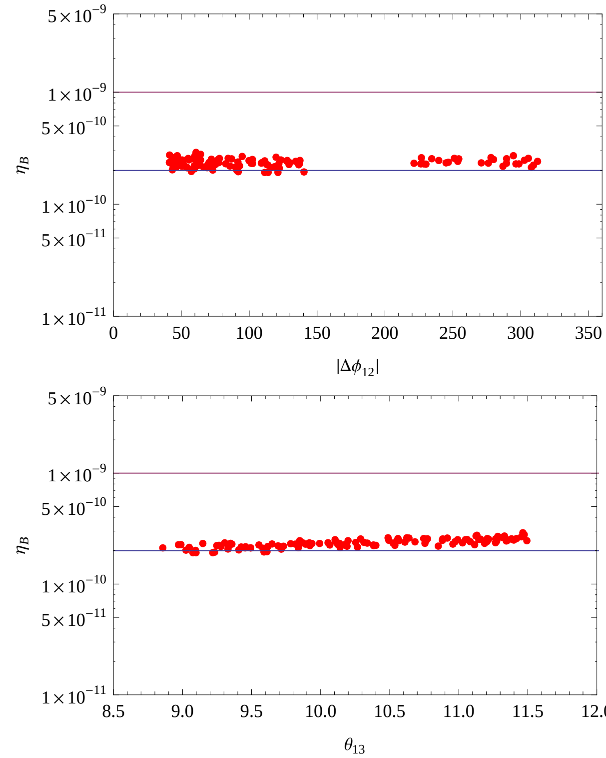


Figure 5: The prediction of flavor-independent leptogenesis over the phase difference $\Delta\phi_{12}$ [upper panel] and the predicted of the mixing angle $\theta_{13}$ [lower panel] with $\tan\beta = 45$. The horizontal lines are the current bound of $\eta_B$..

Below temperatures $T \sim M_i \lesssim (1+\tan^2\beta)\times 10^9$ GeV muon and tau charged lepton Yukawa interactions are much faster than the expansion $H$ rendering the $\mu$ and $\tau$ Yukawa couplings in equilibrium. Then, in this case the final baryon asymmetry is given by [10] as

$$Y_B \simeq \frac{10}{31}\sum_{i=1}^{2}\Big[Y_i^e\Big(\varepsilon_i^e, \frac{93}{110}\widetilde{m}_i^e\Big) + Y_i^\mu\Big(\varepsilon_i^\mu, \frac{19}{30}\widetilde{m}_i^\mu\Big) + Y_i^\tau\Big(\varepsilon_i^\tau, \frac{19}{30}\widetilde{m}_1^\tau\Big)\Big] \quad (30)$$

Finally, the baryon-to-photon is obtained as in the flavor independent leptogenesis. From Eqs. (21) and (28), it is easy to realize that the flavored CP-asymmetries get enhanced over that of the un-flavored ones $\varepsilon_i^\alpha/\varepsilon_i \sim 1/y_\tau^2 t$ due to flavor effects.

## 5. Numerical analysis and summary

Confronting neutrino masses and mixing in the context of our scheme with low energy data given in Eq. (7), we determine the allowed regions of the model parameters for which we estimate the lepton asymmetry. In the Figs. 4 and 5, we plot the predictions of baryon asymmetries $\eta_B$ as functions of the phase difference $\Delta\phi_{12}$ (lower panel) and the mixing angle $\theta_{13}$ (upper panel) imposed initially at GUT scale for $M = 5\times 10^6$ GeV, $M_3 = 10^{12}$ GeV. The Fig. 4 corresponds to the flavor-dependent leptogenesis with $\tan\beta = 10$; where as the flavor-independent leptogenesis with $\tan\beta = 45$ is shown in the Fig. 5. The horizontal lines are the bound of $\eta_B$ given in Ref. [11].

$$2\times 10^{-10} \le \eta_B \le 10\times 10^{-10} \quad (68\% CL). \qquad (31)$$

It is interesting to see that, the generated baryon asymmetry in our scenario is almost constant (for each value of $\tan\beta$). The allowed value of the mixing angle $\theta_{13}$ is rather large $8.6^0 \le \theta_{13} \le 11.6^0$. Besides, the flavor-independent baryon asymmetry value is rather narrow, and that of flavor-dependent case is much wider. The reason is that, the flavor-dependent leptogenesis, as shown in Eq. (30), has the contribution of many sources which due to flavor effects.

As explained in Sec. 4.1, the flavored leptogenesis enhances over the conventional leptogenesis at the same value of $\tan\beta$. It is worthy to note that, in our work, the flavor-independent leptogenesis is successful generated with appropriate values of $\tan\beta$ which is absented as predicted by the Ref. [3].

As a summary, we have considered an exact MTRS in neutrino sector realized at GUT scale in the context of the seesaw model. The exact MTRS, which is realized in the basis where charged lepton and heavy neutrino mass matrices are diagonal, leads to vanishing lepton asymmetries. We have shown that, in the MSSM seesaw model, the RG evolution from the GUT scale to the seesaw scale can induce a successful letogenesis. The right amount of of the baryon asymmetries $\eta_B$ have been achieved via so-called resonant leptogenesis, both flavor independent and flavor dependent cases. We have found that the flavored leptogenesis is enhanced over the unflavored one. In our scenario, the seesaw scale is rather small, so the well-known gravitino problem is safely avoided.

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