Leptogenesis in a Left-Right Symmetric Model with double seesaw

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Abstract. We explore the connection between the low-scale CP-violating Dirac phase (δ) and high-scale leptogenesis in a Left-Right Symmetric Model (LRSM) with scalar bidoublet and doublets. The model's fermion sector includes one sterile neutrino (S_L) per generation to enable a double seesaw mechanism in the neutral fermion mass matrix, implemented by performing type-I seesaw twice. The first seesaw generates the Majorana mass term for heavy right-handed (RH) neutrinos (N_R), and in the second, the light neutrino mass is linearly dependent on S_L mass. We use charge conjugation (C) as the discrete left-right (LR) symmetry, aiding in deriving the Dirac neutrino mass matrix (M_D) in terms of light and heavy RH neutrino masses and the light neutrino mixing matrix U_{PMNS} (containing δ). We demonstrate the feasibility of unflavored thermal leptogenesis via RH neutrino decay using the obtained M_D and RH neutrino masses as input. A thorough analysis of the Boltzmann equations describing asymmetry evolution is conducted in the unflavored regime, showing that the CP-violating Dirac phase alone can generate the required leptonic asymmetry for given input parameters, with or without Majorana phases. Finally, we discuss constraining our model with current and upcoming oscillation experiments aimed at refining the value of δ.

1 Introduction

The LRSM [2, 3] emerged as a theoretical expansion of the Standard Model (SM) of particle physics, aiming to treat left- and right-handed fermion fields equivalently. This framework naturally elucidates observed parity violations in weak interactions through the dynamics of symmetry breaking [4]. LRSM encompasses the entire range of chiral fermions, including right-handed neutrinos (N_R), which can possess both Dirac and Majorana mass terms contingent upon the scalar content of the model. The choice of scalars also governs the model's phenomenological aspects.

Leptogenesis is a well-studied phenomenon in particle physics that seeks to explain the observed asymmetry in the matter-antimatter content of the Universe, a significant discrepancy in the SM. The lepton/antilepton number density asymmetry in leptogenesis is established before the electroweak phase transition. This asymmetry is subsequently converted

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into the observed baryon asymmetry of the Universe (BAU) through higher-dimensional anomalous $B + L$ violating sphaleron processes [5]. One of the early implementations of leptogenesis, proposed by Fukugita and Yanagida [6], involves CP-violating and lepton number-violating out-of-equilibrium decay of heavy neutrinos introduced as singlets under the SM gauge group. This type of leptogenesis is referred to as thermal leptogenesis, where the masses of right-handed neutrinos are typically assumed to be very high, around 10^{10} GeV or higher.

Initially, in the context of LRSM with doublet scalars, no Majorana mass terms were feasible for heavy right-handed neutrinos, prohibiting leptogenesis. An extension of the fermion sector with three generations of sterile neutrinos (S_L) enables a double type-I seesaw [7] in the neutrino mass matrix, thereby imparting a Majorana nature to right-handed neutrinos. This framework has been applied to examine neutrinoless double beta decay [8]. We employ a similar framework to explore thermal unflavored leptogenesis and establish a direct link between low- and high-energy CP violations. In our numerical analysis, we focus on the dependence of BAU on the low-scale CP-violating phase (δ). The coincidence of light and heavy neutrino mixings in our framework allows the dynamical evolution of baryon sector asymmetry to depend on the light neutrino sector. This connection not only reduces the number of input parameters but also strengthens the viability of our model.

2 Model framework and motivation

In the LRSM, the SM gauge group is extended to

$$\mathcal{G}_{LR} \equiv SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \tag{1}$$

where $B - L$ represents the difference between baryon (B) and lepton (L) numbers. The electric charge Q is defined as

$$Q = T_{3L} + T_{3R} + \frac{B-L}{2}. \tag{2}$$

Here, T_{3L} and T_{3R} are, respectively, the third components of isospin of the gauge groups $SU(2)_L$ and $SU(2)_R$. The fermion spectrum of this model comprises all the SM fermions plus a right-handed neutrino N_R. The fermion fields with their quantum numbers ($SU(2)_L \times SU(2)_R \times U(1)_{B-L}$) are as follows:

$$q_L = \begin{pmatrix} u_L \\ d_L \end{pmatrix} \equiv [2, 1, 1/3]\,, \; q_R = \begin{pmatrix} u_R \\ d_R \end{pmatrix} \equiv [1, 2, 1/3]\,,$$
$$\ell_L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix} \equiv [2, 1, -1]\,, \quad \ell_R = \begin{pmatrix} N_R \\ e_R \end{pmatrix} \equiv [1, 2, -1]\,.$$

We have dropped the $SU(3)_C$ quantum numbers for simplicity. The model's scalar sector comprises Higgs bidoublet Φ and the Higgs doublets: H_L and H_R. The matrix structures of the scalar fields are,

$$\Phi = \begin{pmatrix} \phi_1^0 & \phi_2^+ \\ \phi_1^- & \phi_2^0 \end{pmatrix} \sim [2, 2, 0], \;\; H_L = \begin{pmatrix} h_L^+ \\ h_L^0 \end{pmatrix} \sim [2, 1, 1], \;\; H_R = \begin{pmatrix} h_R^+ \\ h_R^0 \end{pmatrix} \sim [1, 2, 1].$$

The spontaneous symmetry breaking (SSB) scheme from LRSM to SM to $U(1)_{em}$ is as follows:

SSB of LRSM

$$SU(2)_L \times \boxed{SU(2)_R \times U(1)_{B-L}} \xrightarrow{\langle H_R(1,2,1)\rangle} SU(2)_L \times U(1)_Y \xrightarrow[\subset\Phi(2,2,0)]{\langle\phi(1_L,1/2_Y)\rangle} U(1)_{em} \tag{3}$$

The right-handed gauge bosons $W_R^\pm$ and Z' get their masses from the VEV $\langle H_R^0\rangle \equiv v_R$ after SSB of LRSM to SM. H_L does not participate in SSB, but it is required in the particle spectrum for left-right invariance. The electroweak symmetry breaking ($SU(2)_L \times U(1)_Y \to U(1)_{em}$) is achieved by assigning non-zero VEVs: $\langle\phi_1^0\rangle \equiv v_1$ and $\langle\phi_2^0\rangle \equiv v_2$ to the neutral components of Higgs bidoublet Φ, with $v = \sqrt{v_1^2 + v_2^2} \simeq 246$ GeV. The Yukawa Lagrangian with usual quarks and leptons reads as,

$$-\mathcal{L}_{Yuk} \supset \overline{q_L}\left[Y_1\Phi + Y_2\widetilde{\Phi}\right]q_R + \overline{\ell_L}\left[Y_3\Phi + Y_4\widetilde{\Phi}\right]\ell_R + \text{h.c.}\,. \tag{4}$$

Here $\widetilde{\Phi} = \sigma_2\Phi^*\sigma_2$ and σ_2 is the second Pauli matrix. When the neutral components (ϕ_1^0 and ϕ_2^0) of Higgs bidoublet acquire non-zero VEVs and give masses to quarks, charged leptons and light neutrinos (Dirac mass) as follows,

$$\begin{aligned} M_u &= Y_1 v_1 + Y_2 v_2\,, \qquad M_d = Y_1 v_2 + Y_2 v_1\,,\\ M_e &= Y_3 v_2 + Y_4 v_1\,, \qquad M_D^\nu \equiv M_D = Y_3 v_1 + Y_4 v_2\,. \end{aligned} \tag{5}$$

Here, M_u and M_d are the up-type and down-type quark mass matrices, and M_e and M_D represent charged lepton and Dirac neutrino mass matrices, respectively. In contrast to Yukawa couplings, which are complex for non-zero CP-asymmetry [9], VEVs v_1 and v_2 are here assumed to be real. If we have $v_2 \ll v_1$ and $|Y_3| \ll |Y_4|$, this will result in small Dirac neutrino masses. With these assumptions, the charged lepton and light neutrino masses can be re-expressed as:

$$M_e \approx Y_4 v_1\,, \qquad M_D = v_1\left(Y_3 + M_e\frac{v_2}{v_1^2}\right) \approx vY_3 \equiv vY_D\,. \tag{6}$$

Here, $v = \sqrt{v_1^2 + v_2^2} \approx v_1$ for $v_2 \ll v_1$. For the left and right gauge coupling constants to be equal, i.e. $g_L = g_R$, we need an additional discrete LR symmetry. The choice of this LR symmetry is twofold [10]: (i.) a generalized parity $\mathcal{P}$ and (ii.) a generalized charge conjugation C. Under the parity and charge conjugation operations, the fields transform as follows:

$$\mathcal{P}: \begin{cases} \ell_L \leftrightarrow \ell_R, \quad q_L \leftrightarrow q_R, \\ \Phi \leftrightarrow \Phi^\dagger, \quad H_L \leftrightarrow H_R, \quad \tilde{\Phi} \leftrightarrow \tilde{\Phi}^\dagger \end{cases} \Bigg| \quad C: \begin{cases} \ell_L \leftrightarrow \ell_R^c, \quad q_L \leftrightarrow q_R^c, \\ \Phi \leftrightarrow \Phi^T, \quad H_L \leftrightarrow H_R^*, \quad \tilde{\Phi} \leftrightarrow \tilde{\Phi}^T. \end{cases} \tag{7}$$

Imposition of either of these discrete symmetries in LRSM makes the Lagrangian in eq. (4) invariant, and it leads to Hermitian Yukawa matrices for the case of discrete $\mathcal{P}$ symmetry. For the case of discrete C symmetry, the Yukawa matrices become symmetric. In our discussion, we consider C symmetry as the additional discrete symmetry.

The Lagrangian in eq. (4) has no lepton number or equivalently (through sphaleron process) no baryon number violating term, one of the three Sakharov conditions. Thus, to have successful leptogenesis in the above-mentioned framework, which is the aim of this work, we require to extend the fermion sector with one additional fermion gauge singlet $S_L \sim (1,1,0)$ ($S_L \overset{C}{\leftrightarrow} S_L^C$) per generation. The addition of fermion gauge singlets S_L ensures lepton number violation (LNV) by induced Majorana mass terms through the double seesaw mechanism.

2.1 Double Seesaw and Neutrino Masses

As discussed in [8], addition of sterile neutrinos (fermion gauge singlets) S_L enable to implement the double seesaw mechanism within the manifest LRSM. The relevant interaction Lagrangian for generation of fermion masses is given by:

$$\begin{aligned}\mathcal{L}_{LRDSM} &= -\mathcal{L}_{M_D} - \mathcal{L}_{M_{RS}} - \mathcal{L}_{M_S} \\ &= -\sum_{\alpha,\beta} \overline{\nu_{\alpha L}}[M_D]_{\alpha\beta} N_{\beta R} - \sum_{\alpha,\beta} \overline{S_{\alpha L}}[M_{RS}]_{\alpha\beta} N_{\beta R} - \frac{1}{2}\sum_{\alpha,\beta} \overline{S^c_{\alpha R}}[M_S]_{\alpha\beta} S_{\beta L} + \text{h.c.}\end{aligned} \tag{8}$$

Here $\mathcal{L}_{M_D}$ and $\mathcal{L}_{M_{RS}}$ are the Dirac mass terms connecting $\nu_L - N_R$ and $N_R - S_L$ respectively. The term $\mathcal{L}_{M_S}$ represents the bare Majorana mass term for sterile neutrinos S_L. In eq. (8), $S^c_{\alpha R} \equiv C(\overline{S_{\alpha L}})^T$, C stands for charge conjugation operation ($C = i\gamma^2\gamma^0$). We note that the Higgs doublet H_L in our model framework is required just for left-right invariance, and it does not participate in SSB. Hence $< H_L^0 >= 0$ and it prevents the mass term connecting $\nu_L - S_R^c$ through the interaction $\sum_{\alpha,\beta} \overline{\ell_{\alpha L}}(Y_{LS})_{\alpha\beta}\widetilde{H_L} S^c_{\beta R}$+ h.c..

After the scalar fields acquire VEVs and thus lead to SSB, the total 9×9 neutral fermion mass matrix in the flavor basis ($\left(\nu_L, N_R^c, S_L\right)$) becomes

$$\mathcal{M}_{LRDSM} = \begin{bmatrix} \mathbf{0} & M_D & \mathbf{0} \\ M_D^T & \mathbf{0} & M_{RS} \\ \mathbf{0} & M_{RS}^T & M_S \end{bmatrix} \tag{9}$$

With the assumed hierarchy $|M_D| \ll |M_{RS}| < |M_s|$, we apply double seesaw approximate block diagonalization to $\mathcal{M}_{LRDSM}$. This gives us the mass matrices of light and heavy neutrinos as follows:

$$\begin{aligned} m_\nu &\cong M_D (M_{RS}^T)^{-1} M_S M_D^T M_{RS}^{-1} \text{ (mass matrix of light neutrinos)}, \\ m_N &\equiv M_R \cong -M_{RS} M_S^{-1} M_{RS}^T \text{ (mass matrix of RH neutrinos)}, \\ m_S &\cong M_S \text{ (mass matrix of sterile neutrinos)}. \end{aligned} \tag{10}$$

The detailed discussion on double seesaw approximations and derivations of mass matrices expressed in eq. (10) can be referred from the ref. [8].

3 Connecting low-scale and high-scale CP violation

It is important to establish a connection between the parameters at low-energy (particularly the CP-violating phases), which can be probed in present and future experiments, and at high-energy that would be relevant for leptogenesis. For an unflavored analysis (i.e. considering only N_1 decays), this CP-asymmetry (ϵ_1) is defined as

$$\epsilon_1 = \frac{\Gamma(N_1 \to \ell\Phi) - \Gamma(N_1 \to \ell^c\Phi^c)}{\Gamma_D}. \tag{11}$$

Here, Γ_D is the total decay rate for N_1 and is expressed as $\frac{[Y_D^\dagger Y_D]_{11} m_{N_1}}{8\pi}$, where Y_D is Yukawa coupling matrix as introduced in eq. (6). The asymmetry ϵ_1 arises from the interference between tree and 1-loop wave and vertex diagrams as shown in Fig. 1. We work under the assumption that the RH neutrinos have the mass hierarchy as $m_{N_1} < m_{N_2} < m_{N_3}$, so that it is the decays of the N_1 that essentially determines the sought-after asymmetry [9].

$$\epsilon_1 \approx -\frac{3.m_{N_1}}{16\pi v^2 (M_D^\dagger M_D)_{11}} \left[\frac{Im[(M_D^\dagger M_D)^2_{21}]}{m_{N_2}} + \frac{Im[(M_D^\dagger M_D)^2_{31}]}{m_{N_3}}\right]. \tag{12}$$

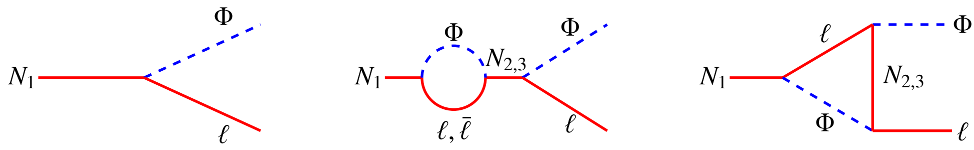

Figure 1: Tree, 1-loop and vertex diagrams for heavy neutrino decay. The asymmetry ϵ_1 results from interference of the 1-loop diagrams with tree level coupling.

Here, $M_D = v.Y_D$. From eq. (12), we see that to determine the CP-asymmetry, we need to find the structure of M_D and the masses of right-handed neutrinos (m_{N_i}). The CI parametrization is widely used in which Y_D is expressed in terms of a complex orthogonal matrix R [11]. In our approach, the theoretical framework of double seesaw in LRSM with considered particle spectrum incorporated with discrete LR symmetry C and the screening condition naturally leads to a direct connection between low- and high-energy CP violations.

3.1 Screening effect

For the structure of the light neutrino mass matrix, m_ν, to be determined by the structure of the sterile neutrino mass matrix, m_S, we apply screening (cancellation) of Dirac structures in the expression of m_ν of eq. (10) [12]. In ref. [8], the authors achieve it by considering M_D and M_{RS} to be proportional to identity (I). The consideration was relevant for studying neutrinoless double beta decay. However, we cannot have this screening condition for studying leptogenesis as it leads to vanishing CP-asymmetry ($\epsilon_1 = 0$). We take the screening condition [13]:

$$M_D = \frac{1}{k} M_{RS}^T \tag{13}$$

Here k is a real constant. It is worth noting that the renormalization group (RG) running of Yukawa couplings (Y_D and Y_{RS}) from the LRSM symmetry breaking scale to the EW scale might affect the screening mechanism. While ref. [13] mentions that RG running would not destroy the screening, one may refer to ref. [12] for a detailed discussion on the impact of RGE effect on the screening in the context of seesaw mechanism and neutrino masses. With eq. (13), the relation between light neutrino and sterile neutrino mass matrices m_ν and m_S becomes $m_S = k^2 m_\nu$. The light neutrino Majorana mass matrix is diagonalized with the Pontecorvo, Maki, Nakagawa, Sakata (PMNS) mixing matrix $U_{PMNS} \equiv U_\nu$:

$$\hat{m}_\nu = U_\nu^\dagger m_\nu U_\nu^* = \mathrm{diag}(m_1, m_2, m_3),\ m_i > 0 \text{ for } i = 1, 2, 3. \tag{14}$$

In what follows, we will use the standard parametrization of the PMNS matrix [14]:

$$U_{\rm PMNS} \equiv U_\nu = \begin{pmatrix} c_{13}c_{12} & c_{13}s_{12} & s_{13}e^{-i\delta} \\ -c_{23}s_{12} - c_{12}s_{13}s_{23}e^{i\delta} & c_{12}c_{23} - s_{12}s_{13}s_{23}e^{i\delta} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}s_{13}c_{23}e^{i\delta} & c_{13}c_{23} \end{pmatrix} \underbrace{\begin{pmatrix} 1 & 0 & 0 \\ 0 & e^{i\alpha/2} & 0 \\ 0 & 0 & e^{i\beta/2} \end{pmatrix}}_{\text{Majorana phase matrix}} \tag{15}$$

where δ is the Dirac CP phase ($0 \le \delta \le 2\pi$) and α, β are the Majorana CP phases ($0 \le \alpha, \beta \le 2\pi$). All the other parameters have their usual meanings. The sterile neutrino Majorana mass matrix m_S is diagonalized by a unitary matrix U_S as $\hat{m}_S = U_S^\dagger m_S U_S^*$, where we have

$\hat{m}_S = \text{diag}(m_{S_1}, m_{S_2}, m_{S_3})$, $m_{S_k} > 0$, $k = 1, 2, 3$. Since $m_S = k^2 m_\nu$, the diagonalization can be done with the same mixing matrix U_ν, i.e. $U_S = U_\nu$.
So for the considered scenario, the light neutrino masses m_i and the sterile neutrino masses m_{S_k} are related as

$$m_i = \frac{1}{k^2} m_{S_i},\ i = 1, 2, 3 \tag{16}$$

For normal ordering (NO) mass spectrum of active neutrinos ($m_1 < m_2 < m_3$), we have:

$$m_1 = \text{ lightest neutrino mass} \quad m_2 = \sqrt{m_1^2 + \Delta m_{\rm sol}^2}, \quad m_3 = \sqrt{m_1^2 + \Delta m_{\rm atm}^2}. \tag{17}$$

Therefore from eq. (16), the relations between masses of active and sterile neutrinos become:

$$m_{S_1} = \frac{m_1}{m_3} m_{S_3}, \quad m_{S_2} = \frac{m_2}{m_3} m_{S_3}, \quad m_{S_1} < m_{S_2} < m_{S_3}. \tag{18}$$

For inverted ordering (IO) ($m_3 < m_1 < m_2$), we have:

$$m_3 = \text{ lightest neutrino mass} \quad m_1 = \sqrt{m_3^2 + \Delta m_{\rm atm}^2}, \quad m_2 = \sqrt{m_3^2 + \Delta m_{\rm sol}^2 + \Delta m_{\rm atm}^2}, \tag{19}$$

and masses relations from eq. (16) become:

$$m_{S_1} = \frac{m_1}{m_2} m_{S_2}, \quad m_{S_3} = \frac{m_3}{m_2} m_{S_2}, \quad m_{S_3} < m_{S_1} < m_{S_2}. \tag{20}$$

In both orderings, we have $\Delta m_{\rm sol}^2 = \Delta m_{21}^2$ and $\Delta m_{\rm atm}^2 = \left|\Delta m_{31}^2\right|$ [15].

3.2 Choice of basis

We work in the basis where the charged lepton mass matrix is diagonal. The right-handed neutrino Majorana mass matrix m_N can be diagonalized by a unitary matrix U_N as $\hat{m}_N = U_N^\dagger m_N U_N^*$. Here, we have $\hat{m}_N = diag(m_{N_1}, m_{N_2}, m_{N_3})$ with $m_{N_i} (i = 1, 2, 3)$ being the mass of the heavy RH Majorana neutrino N_i. By using the screening result ($U_S = U_\nu$), we have

$$\hat{m}_S = U_\nu^\dagger m_S U_\nu^* \qquad \Longrightarrow \quad m_S^{-1} = U_\nu^* \hat{m}_S^{-1} U_\nu^\dagger. \tag{21}$$

Using m_S^{-1} from eq. (21) in the expression of m_N from eq. (10), we have:

$$\hat{m}_N = -\underbrace{U_N^\dagger M_{RS} U_\nu^*}\, \hat{m}_S^{-1}\, \underbrace{U_\nu^\dagger M_{RS}^T U_N^*} \tag{22}$$

For eq. (22) to be consistent, the right-hand side should be diagonal. As $\hat{m}_S^{-1}$ is diagonal, so

$$U_N^\dagger M_{RS} U_\nu^* = \hat{m}_{RS}. \tag{23}$$

Since we have considered C symmetry as the additional discrete symmetry in our model framework, we have M_D and through eq. (13), M_{RS} as the symmetric matrices[1]. Thus from eq. (23), we must have

$$U_N = U_\nu. \tag{24}$$

Therefore, the diagonalization of RHN Majorana mass matrix m_N can be performed with the same mixing matrix U_ν. The screening condition eq. (13) also modifies to:

$$M_D = \frac{1}{k} M_{RS} \tag{25}$$

[1]This is one of the reasons for considering C symmetry as the additional discrete LR symmetry in our model framework.

3.3 Determining M_D

We derive here the matrix structure of M_D analytically and show the connection between low- and high-energy CP violations. From eq. (10), the heavy neutrino mass matrix is given as, $m_N = -M_{RS} m_S^{-1} M_{RS}^T$. By using $m_S = k^2 m_\nu$ and noting that M_D, M_{RS} are symmetric matrices, we can rewrite it as

$$m_N = -\frac{1}{k^2} M_{RS} m_\nu^{-1} M_{RS}. \tag{26}$$

Simplifying for M_{RS}, we get the final expression as:

$$M_{RS} = m_\nu \sqrt{-k^2 m_\nu^{-1} m_N}. \tag{27}$$

Now by using the equations (24) and (25), and extracting the square root of matrices in eq. (27), we get the expression for M_D:

$$\boxed{M_D = \frac{1}{k} M_{RS} = i.U_\nu \hat{m}_\nu (\hat{m}_\nu^{-1} \hat{m}_N)^{1/2} U_\nu^T} \tag{28}$$

For the full derivation of M_D, one may refer the appendix in [1]. We have thus connected low- and high-scale CP violations through eq. (28), as M_D is expressed in terms of $U_\nu \equiv U_{PMNS}$. The masses of heavy neutrinos are denoted by m_{N_i} $(i = 1, 2, 3)$, with $m_{N_1} < m_{N_2} < m_{N_3}$.

4 Leptogenesis

The scenario of leptogenesis to produce an asymmetry requires fulfillment of three Sakharov conditions [16] as the minimum necessary criteria. Thus, in this section, we present an implementation of thermal leptogenesis in the context of our framework by ensuring the sanctity of the required Sakharov criteria and also by producing the adequate lepton asymmetry required for meeting the observational evidence of BAU from the Planck [17] data. The complex nature of Yukawa coupling matrices (Y_D) ensures that the decays are CP-violating. To achieve out-of-equlibrium conditions, the mass of N_R is taken to be large so that the temperature of the Universe at the time of their decay is less than their rest mass, and hence the probability of inverse interactions decreases, keeping the decays out-of thermal equilibrium. To highlight the connection between this asymmetry parameter, ϵ_1 (relevant for leptogenesis) and the CP-violating phase, δ (coming from low-scale physics), we derive the various terms in eq. (12) analytically in terms of the CP-violating phase, using the structure of M_D given in eq. (28) along with the values of experimental parameters given in ref. [15] and by fixing rest of the input parameters.

4.1 Normal Ordering

From eq. (12), for normal ordering (NO) mass spectrum of active neutrinos, we have:

$$\begin{aligned}
\mathrm{Im}[(M_D^\dagger M_D)_{21}^2] &= \sin\delta\,(2.06 + 16\cos\delta) \times 10^6\ \mathrm{GeV}^2 \\
\mathrm{Im}[(M_D^\dagger M_D)_{31}^2] &= \sin\delta(-2.06 + 11.86\cos\delta) \times 10^6\ \mathrm{GeV}^2 \\
(M_D^\dagger M_D)_{11} &= 1042.64\ \mathrm{GeV}^2
\end{aligned} \tag{29}$$

Here, we have used the benchmark point values $m_{N_1} = 10^{13}$ GeV, $m_{N_2} = 1 \times 10^{14}$ GeV, $m_{N_3} = 5 \times 10^{14}$ GeV and $m_1 = 0.01$ eV. In the left panel of figure 2, we plot the variation of ϵ_1 against the allowed range of CP-violating phase δ $[0, 2\pi]$. Using the benchmark values of all the input parameters along with the best-fit value of $\delta = 1.08\pi$ [15], and the standard Higgs VEV of $v \simeq 246$ GeV, the asymmetry parameter here is numerically obtained as,

$$\epsilon_1 \sim -3.8 \times 10^{-4}. \tag{30}$$

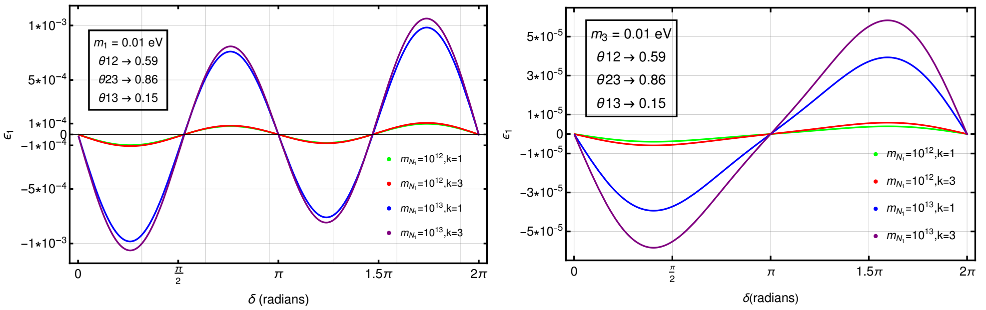

Figure 2: Plot for the dependence of ϵ_1 on CP-violating Dirac phase, δ for the **NO case** in the left panel and for the **IO case** in the right panel for different combinations of right-handed neutrino mass and the hierarchy in that sector. Value of variable k represents the masses of heavier right-handed neutrino for the structure: $M_{N_2} = 1\times10^k\times M_{N_1}$ and $M_{N_3} = 5\times10^k\times M_{N_1}$.

4.2 Inverted Ordering

Similar to the analysis performed in subsection 4.1, we here present the results for the case of inverted ordering. Thus, from eq. (12), we get:

$$\begin{aligned}
\mathrm{Im}[(M_D^\dagger M_D)_{21}^2] &= \sin\delta\,(9.46 + 2.36\cos\delta)\times 10^5\ \mathrm{GeV}^2\\
\mathrm{Im}[(M_D^\dagger M_D)_{31}^2] &= \sin\delta(-9.46 + 1.75\cos\delta)\times 10^5\ \mathrm{GeV}^2\\
(M_D^\dagger M_D)_{11} &= 2003.03\ \mathrm{GeV}^2 \qquad (31)
\end{aligned}$$

Here, we take the benchmark point values $m_{N_1} = 10^{13}$ GeV, $m_{N_2} = 1\times 10^{14}$ GeV, $m_{N_3} = 5\times 10^{14}$ GeV and $m_3 = 0.01$ eV. In the right panel of figure 2, we plot the variation of ϵ_1 against the allowed range of CP-violating phase δ $[0, 2\pi]$. For a fair comparison, keeping the values of benchmark input parameters the same as that for the NO case along with the best-fit value of $\delta = 1.58\pi$ [15], the asymmetry parameter for the IO case is numerically obtained as,

$$\epsilon_1 \sim +3.92\times 10^{-5}. \qquad (32)$$

Inclusion of Majorana Phases

Our work primarily focuses on the Dirac CP phase (δ) as the exclusive source of CP violation for phenomenological discussions. However, we also incorporate the Majorana phases (α, β) introduced in eq. (15) to encompass all aspects of CP violation and to explore the possibility of lowering the scale of the leptogenesis scenario. Thus, we express the various terms in eq. (12) analytically in terms of the Majorana phases for two different values of the Dirac phase: **(1.)** $\delta = 0$ (CP conserving) and **(2.)** $\delta = 1.08\pi$ (best-fit) for both NO and IO cases.

A)- NO case

(1.) For $\delta = 0$.

$$\begin{aligned}
\mathrm{Im}[(M_D^\dagger M_D)_{21}^2] &\simeq -[5.45\sin\beta + 85.81\sin\alpha]\times 10^{-10}\ \mathrm{GeV}^2 + f(\alpha,\beta)\,O(10^{-26})\\
\mathrm{Im}[(M_D^\dagger M_D)_{31}^2] &\simeq [1.43\sin\alpha - 5.72\sin\beta - 0.16\sin(\alpha-\beta) - 2.29\sin(\alpha+\beta)]\\
&\qquad \times 10^{-10}\ \mathrm{GeV}^2 + f(\alpha,\beta)\,O(10^{-26})\\
(M_D^\dagger M_D)_{11} &= [1042.64 - [2.84\cos\alpha - 11.37\cos(\alpha-\beta) + 2.84\cos\beta]\times 10^{-14}]\ \mathrm{GeV}^2 \qquad (33)
\end{aligned}$$

Here, $f(\alpha,\beta)O(10^n)$ corresponds to the sine and cosine functions of Majorana phases (α,β) with an order of magnitude n or less, where n is an integer. Using the benchmark values of all the input parameters and the standard Higgs VEV of $v \simeq 246$ GeV, the asymmetry parameter here is numerically obtained as:

$$\epsilon_1 \sim [(0.59 + 3.46\cos\beta)\sin\alpha + (6.23 - 2.60\cos\alpha)\sin\beta] \times 10^{-20}. \quad (34)$$

(2.) For $\delta = 1.08\pi$.

$$\begin{aligned} \mathrm{Im}[(M_D^\dagger M_D)_{21}^2] &\simeq 3.34 \times 10^6 \text{ GeV}^2 + f(\alpha,\beta)\,O(10^{-9}) \\ \mathrm{Im}[(M_D^\dagger M_D)_{31}^2] &\simeq 3.37 \times 10^6 \text{ GeV}^2 + f(\alpha,\beta)\,O(10^{-9}) \\ (M_D^\dagger M_D)_{11} &\simeq 1042.64 \text{ GeV}^2 + f(\alpha,\beta)\,O(10^{-14}) \end{aligned} \quad (35)$$

With the benchmark values, the asymmetry parameter is numerically obtained as:

$$\epsilon_1 \sim -3.8 \times 10^{-4} + f(\alpha,\beta)\,O(10^{-13}) \quad (36)$$

B)- IO case

(1.) For $\delta = 0$.

$$\begin{aligned} \mathrm{Im}[(M_D^\dagger M_D)_{21}^2] &\simeq -(5.87\sin\alpha + 0.98\sin\beta + 3.91\sin(\alpha-\beta)) \times 10^{-10} \text{ GeV}^2 \\ &\quad + f(\alpha,\beta)\,O(10^{-27}) \\ \mathrm{Im}[(M_D^\dagger M_D)_{31}^2] &\simeq -(1.48\sin\alpha + 0.74\sin\beta - 5.93\sin(\alpha-\beta)) \times 10^{-10} \text{ GeV}^2 \\ &\quad + f(\alpha,\beta)\,O(10^{-26}) \\ (M_D^\dagger M_D)_{11} &= [2003.03 - [1.71\cos\alpha - 1.14\cos(\alpha-\beta)] \times 10^{-13}] \text{ GeV}^2 \end{aligned} \quad (37)$$

With the values of input benchmark parameters, the asymmetry parameter here is numerically obtained as:

$$\epsilon_1 \sim [(3.03 + 0.29\cos\alpha)\sin\alpha + 1.34\sin(\alpha-\beta) + 0.55\sin\beta] \times 10^{-20}. \quad (38)$$

(2.) For $\delta = 1.58\pi$.

$$\begin{aligned} \mathrm{Im}[(M_D^\dagger M_D)_{21}^2] &\simeq -9.73 \times 10^5 \text{ GeV}^2 + f(\alpha,\beta)\,O(10^{-10}) \\ \mathrm{Im}[(M_D^\dagger M_D)_{31}^2] &\simeq 8.74 \times 10^5 \text{ GeV}^2 + f(\alpha,\beta)\,O(10^{-10}) \\ (M_D^\dagger M_D)_{11} &\simeq 2003.03 \text{ GeV}^2 + f(\alpha,\beta)\,O(10^{-13}) \end{aligned} \quad (39)$$

With the values of the input benchmark parameters, the asymmetry parameter is numerically obtained as follows:

$$\epsilon_1 \sim +3.92 \times 10^{-5} + f(\alpha,\beta)\,O(10^{-20}). \quad (40)$$

In figure 3, we depict the variation of the CP asymmetry parameter (ϵ_1) with respect to the Majorana phases α and β, while holding the Dirac phase (δ) fixed at a CP-conserving value ($\delta = 0$) and at the best-fit value ($\delta = 1.08\pi$) for both NO and IO cases. These plots reveal that the dependence of ϵ_1 on the Majorana phases is several orders of magnitude smaller than its dependence on the Dirac CP phase. Consequently, Majorana phases play an insignificant role in generating the required CP asymmetry. Hence, we do not include the Majorana phases for our Boltzmann analysis to be performed in the subsequent section.

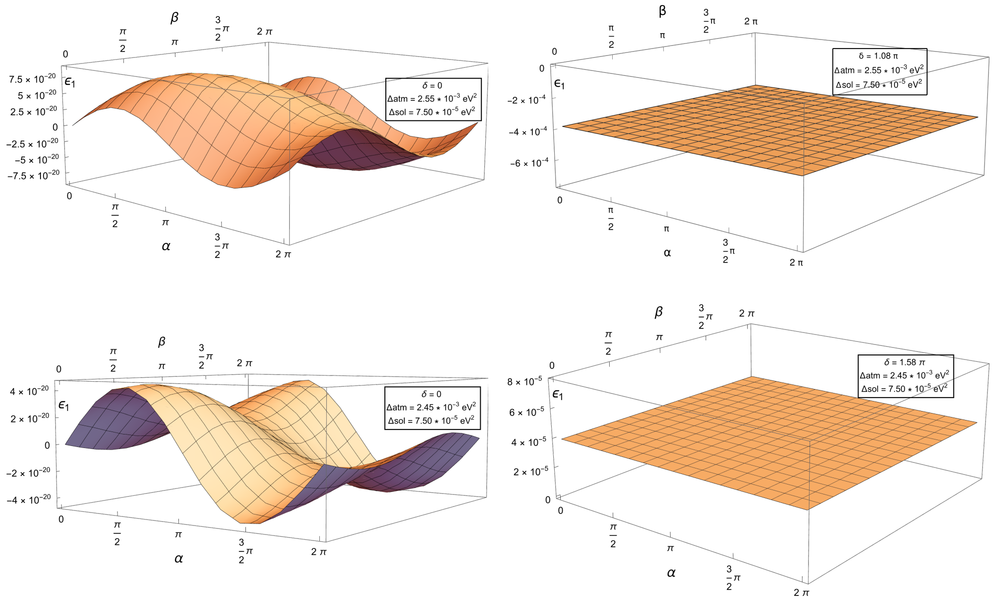

Figure 3: 3D plots depicting the dependence of asymmetry parameter ϵ_1 on both the Majorana phases α and β for the case of zero ($\delta = 0$) and best-fit ($\delta = 1.58\pi$) Dirac phase for the NO case in the top panel and for the IO case in the bottom panel.

4.3 Boltzmann Analysis

Here, we present the structure of flavor singular coupled BEs required for our analysis,

$$\frac{d\eta^{N_1}}{dz} = -\left(\frac{\eta^{N_1}}{\eta^{N}_{\rm eq}} - 1\right)(D_1 + S_1)\,, \tag{41}$$

$$\frac{d\eta^{\Delta L}}{dz} = \varepsilon_1\left(\frac{\eta^{N_1}}{\eta^{N}_{\rm eq}} - 1\right)\widetilde{D}_1 - \frac{2}{3}\eta^{\Delta L}W_l\,, \tag{42}$$

where $z = m_{N_1}/T$ is a dimensionless variable (T being the temperature of the Universe) and $\eta^N_{\rm eq} \equiv n^N_{\rm eq}/n^\gamma = z^2K_2(z)/2\zeta(3)$ is the heavy neutrino equilibrium number density, $K_n(z)$ being the n-th order modified Bessel function of the second kind and $\zeta(3)$ is Riemann zeta function ($\zeta(s)$) evaluated at $s = 3$. n^γ is referred as the comoving photon number density. The various decay (D_1, $\widetilde{D}_1$), scattering (S_1) and washout (W_l) rates appearing in equations (41) and (42) are given by

$$\widetilde{D}_1 = \frac{z}{n^\gamma H_N}\widetilde{\gamma}^D, \quad D_1 = \frac{z}{n^\gamma H_N}\gamma^D, \quad S_1 = \frac{z}{n^\gamma H_N}(\gamma^{S_L} + \gamma^{S_R}), \tag{43}$$

$$W_l = \frac{z}{n^\gamma H_N}\left[\gamma^D + \widetilde{\gamma}^{S_L} + \widetilde{\gamma}^{S_R} + \gamma^{(\Delta L=0)} + \gamma^{(\Delta L=1)} + \gamma^{(\Delta L=2)}\right], \tag{44}$$

where $H_N \equiv H(z = 1) \simeq 17m^2_{N_1}/M_{\rm Pl}$ is the Hubble parameter at $z = 1$, assuming only SM degrees of freedom in the thermal bath, $M_{\rm Pl} = 1.2{\times}10^{19}$ GeV is the Planck mass. The solution to these BEs considering all the decays, inverse-decays, scatterings and washout terms of N_1

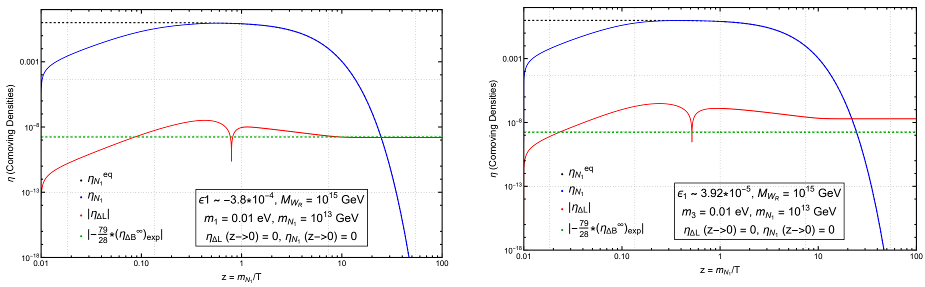

Figure 4: Cosmological evolution of the number densities present in the coupled BEs (41) and (42) for the **NO** case in the left panel and for the **IO** case in the right panel.

for both the NO and IO cases is presented in figure 4. The overall interplay and dynamics of these terms provide an acceptable magnitude for the final asymmetry in the lepton sector close to ($\eta^{\infty}_{\Delta L} \sim -6 \times 10^{-10}$) as shown in the figure. Then, we assume a successful transfer of this asymmetry to the baryonic sector via sphalerons. As the value of asymmetry parameter, (ϵ_1) comes to be negative in this NO case (eq. 30), the sign of final baryon asymmetry, ($\eta^{\infty}_{\Delta B}$) is obtained to be positive here, in accordance with the leptogenesis requirements. But the value of asymmetry parameter, (ϵ_1) for the IO case comes to be positive (eq. 32), hence the sign of final baryon asymmetry, ($\eta^{\infty}_{\Delta B}$) for the given parameter space is obtained to be negative here, thus IO case is ruled out for the chosen parameter space.

5 Conclusion

We have explored thermal unflavored leptogenesis within a category of Left-Right Symmetric Models featuring a scalar bidoublet and doublets, along with the addition of a single copy of sterile neutrino, S_L, per generation in the fermion sector. The requisite source of CP violation for successful leptogenesis is the generic Dirac neutrino mass matrix, (M_D), which establishes a connection between ν_L and N_R. The distinctive aspect of our investigation lies in M_D becoming contingent on the low-energy CP-violating Dirac phase, δ (embedded in U_{PMNS}), without resorting to any parameterization. This presents an intriguing incentive for ongoing long baseline experiments like NOvA, T2K, DUNE, T2HK, T2HKK, and future endeavours like JUNO to explore leptogenesis through δ indirectly.

For thoroughness, we conduct our analysis for both the normal and inverted mass orderings of the light neutrinos. Employing specified benchmark points for the normal ordering case, we obtain a final baryon asymmetry value consistent with the observed $\eta^{\infty}_{\Delta B} = (6.105^{+0.086}_{-0.081}) \times 10^{-10}$. Correspondingly, for similar benchmark values, the final asymmetry achieved in the inverted ordering case exhibits approximately the same order of magnitude, albeit with a negative sign. Remarkably, our analysis reveals that within the considered model framework of double seesaw, the asymmetry parameter (ϵ_1) displays minimal dependency on the Majorana phases α and β for the provided set of input parameters in both the normal and inverted ordering cases. This underscores δ as the principal source for generating the requisite baryon asymmetry. However, with other choices of input parameters, one may observe a distinct dependency of ϵ_1 on α and β, albeit such selections might steer us away from the thermal unflavored regime. We intend to extend this study to examine the impact of non-zero Majorana phases in the flavored or resonant regime of leptogenesis.

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