A MINIMAL FLAVOR MODEL FOR NEUTRINO MASS AND LEPTOGENESIS*

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To explain the observed neutrino mixing, we exploit an A_4 discrete flavor symmetric model, where neutrino masses are generated via type-I seesaw. Thanks to the flavor structure of the model, only the normal hierarchy of neutrino mass is allowed, atmospheric mixing falls in the lower octant, and leptonic CP phases including the Majorana phases also get constrained. Owing to the symmetry, the right-handed neutrinos are degenerate to start with. Once we include the renormalization group running this degeneracy can be lifted and adequate lepton asymmetry can be generated.

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1. Introduction

Over the last few decades, neutrino physics experiments have entered the precision era. However, there are a few issues that still need to be settled which include octant of the atmospheric mixing angle (θ_{23}), magnitude of Dirac CP phase (δ), hierarchy of light-neutrino masses *etc.* In a wider perspective, the lepton mixing is found to be quite different compared to the observed quark mixing. To obtain a deeper understanding of it, one needs to investigate the origin of the neutrino masses and mixing from a symmetry perspective. The type-I seesaw mechanism is the simplest extension of Standard Model (SM) to explain tiny neutrino mass where three singlet right-handed neutrinos (RHN) are included in the most general case. In addition to this, the involvement of flavor symmetries here is surely an interesting possibility to explain the typical mixing pattern in the lepton sector. Several non-Abelian discrete groups have been extensively used in this regard, see reviews [1–3] and references therein. Among many other groups

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used for this purpose, A_4 turns out to be the most economical one which contains a three-dimensional representation (3) unifying the three SM lepton doublets [4, 5]. For example, such conjugation of A_4 discrete symmetry and type-I seesaw were earlier proposed to explain the tri-bimaximal (TBM) lepton mixing pattern [6] in the presence of SM singlet (charged under A_4) flavon fields. However, after precise measurement [7] of the reactor mixing angle θ_{13}, a modification to such models is essential and often it is achieved in a non-minimal way [8–14].

In this paper, we work in the original framework of the Altarelli–Feruglio (AF) model [4] considering only three A_4 flavons. Owing to the considered symmetry, the RHN mass matrix is found to be diagonal. Originated from the product of two A_4 triplets, an antisymmetric contribution to the Dirac neutrino mass matrix performs an instrumental role in generating non-zero θ_{13} [15–18] in our model which was overlooked in an earlier attempt [19]. Here, we find θ_{23} to lie in the lower octant and only the normal hierarchy (NH) of light-neutrino masses is allowed. This is indeed a significant prediction which can be probed in the ongoing and future neutrino experiments. Furthermore, the presence of the RHNs in this type-I-seesaw-based scenario gives us the opportunity to study leptogenesis [20]. As the RHNs are exactly degenerate at tree level, lepton asymmetry [21] can be generated via renormalization group effects as shown in [19]. Moreover, our study differs from [19] in terms of flavon content and we perform a thorough analysis of flavored leptogenesis by solving the Boltzmann equations.

This paper is organized as follows. In Section 2, we present the construction model, and exercise the associated phenomenology in Section 3. In Section 4, we study leptogenesis including lepton flavor effects. Finally, in Section 5, we make our final conclusions.

2. The model

In [22], we have adopted a type-I seesaw with three RHNs($N_{\rm R}$) assisted by $A_4 \otimes Z_2 \otimes Z_3$ discrete with involvement of the flavons Φ, Ψ, and φ as considered in the original AF [4]. Here, φ transforms as singlet and $N_{\rm R}$, Φ, Ψ transform as triplet under A_4. The additional $Z_2 \otimes Z_3$ symmetry with A_4 forbids many undesired contributions and leads to specific forms for the leptonic mass matrices. In Table 1, we mention the transformation properties of all SM (and beyond) particles involved in our study.

Following Table 1, allowed effective Lagrangian involving charged leptons can be written as

$$\mathcal{L}_{\rm CL} = \frac{y_1^\ell}{\Lambda}\left(\bar{\ell}\,\Phi\right)_{\mathbf{1}} H\, e_{\rm R} + \frac{y_2^\ell}{\Lambda}\left(\bar{\ell}\,\Phi\right)_{\mathbf{1}''} H\, \mu_{\rm R} + \frac{y_3^\ell}{\Lambda}\left(\bar{\ell}\,\Phi\right)_{\mathbf{1}'} H\, \tau_{\rm R} + \text{h.c.}\,, \quad (1)$$

Table 1. Transformation of the fields under $A_4 \otimes Z_2 \otimes Z_3$ symmetry.

Fields	ℓ	$e_{\rm R}$, $\mu_{\rm R}$, $\tau_{\rm R}$	$N_{\rm R}$	H	φ	Φ	Ψ
A_4	3	1 , $1'$, $1''$	3	1	1	3	3
Z_2	1	1	-1	1	-1	1	-1
Z_3	ω	1	1	1	ω	ω	ω

and, similarly, for neutrinos, the Lagrangian is given by

$$\begin{aligned}\mathcal{L}_\nu &= \frac{y_1^\nu}{\Lambda}\left[\left(\bar{\ell}\, N_{\rm R}\right)_{\mathbf{s}} \Psi\right]_{\mathbf{1}} \tilde{H} + \frac{y_2^\nu}{\Lambda}\left[\left(\bar{\ell}\, N_{\rm R}\right)_{\mathbf{a}} \Psi\right]_{\mathbf{1}} \tilde{H} + \frac{y_3^\nu}{\Lambda}\left(\bar{\ell}\, N_{\rm R}\right)_{\mathbf{1}} \varphi\, \tilde{H} \\ &+\frac{1}{2} M\left(\bar{N}_{\rm R}^c N_{\rm R}\right) + \text{h.c.}\,, \end{aligned} \tag{2}$$

where $y_{i=1,2,3}^{\ell,\nu}$ are the associated coupling constants, Λ stands for cut-off scale, and subscripts $\mathbf{s}$, $\mathbf{a}$ correspond to symmetric and antisymmetric parts of products of A_4 triplets. For a detailed description of A_4 multiplication rules, the readers are referred to [1, 3, 23]. Once flavor symmetry is broken in the direction [13, 15, 19, 24–26]

$$\langle\varphi\rangle = v_\varphi\,, \qquad \langle\Phi\rangle = v_\Phi\,(1,1,1)\,, \qquad \langle\Psi\rangle = v_\Psi\,(0,1,0)\,, \tag{3}$$

from Eq. (1), after electroweak symmetry is also broken, the charged lepton mass matrix can be written as

$$Y^\ell = v \begin{pmatrix} f_1^\ell & f_2^\ell & f_3^\ell \\ f_1^\ell & \omega\, f_2^\ell & \omega^2\, f_3^\ell \\ f_1^\ell & \omega^2\, f_2^\ell & \omega\, f_3^\ell \end{pmatrix}; \qquad f_i^\ell = \frac{v_\Phi}{\Lambda}\, y_i^\ell \quad \text{with} \quad i = 1,2,3\,, \tag{4}$$

where $v = 174$ GeV stands for the vacuum expectation value of the SM Higgs. In a similar way, from Eq. (2), the Lagrangian for neutrino sector yields Dirac Yukawa and Majorana mass matrices as

$$Y^\nu = \begin{pmatrix} f_3^\nu & 0 & f_1^\nu - f_2^\nu \\ 0 & f_3^\nu & 0 \\ f_1^\nu + f_2^\nu & 0 & f_3^\nu \end{pmatrix}, \tag{5}$$

$$M_{\rm R} = \begin{pmatrix} M & 0 & 0 \\ 0 & M & 0 \\ 0 & 0 & M \end{pmatrix}, \tag{6}$$

with $f_i^\nu = \frac{v_\Psi}{\Lambda} y_i^\nu$, $i = 1,2,3$. Here, f_2^ν in Eq. (5) represents the antisymmetric contribution to the Dirac neutrino Yukawa, originated from the products of

A_4 triplets (ℓ, $N_{\rm R}$, and Φ) and plays a crucial role [15–18] in realizing correct neutrino oscillation data. Now, the charged lepton matrix given in Eq. (4) can be diagonalized by a matrix given by

$$V = \frac{1}{\sqrt{3}}\begin{pmatrix} 1 & 1 & 1 \\ 1 & \omega & \omega^2 \\ 1 & \omega^2 & \omega \end{pmatrix}, \tag{7}$$

where ω ($= {\rm e}^{2i\pi/3}$) is the cube root of unity. Via type-I seesaw, the effective light-neutrino mass matrix can be obtained as

$$m_\nu = -m_{\rm D} M_{\rm R}^{-1} m_{\rm D}^T\,. \tag{8}$$

Substituting $m_{\rm D}$ and $M_{\rm R}$, one can obtain

$$m_\nu = -\frac{1}{M} V^\dagger \left(v^2 Y^\nu Y^{\nu T}\right) V^*\,, \tag{9}$$

where

$$m_{\rm D} = v V^\dagger Y^\nu = v \mathcal{Y}^\nu\,. \tag{10}$$

Now, the light-neutrino mass matrix given by Eq. (9) can be diagonalized via

$$m_\nu \equiv U {\rm diag}(m_1, m_2, m_3) U^T\,, \tag{11}$$

where $m_{1,2,3}$ are the real positive mass eigenvalues and the effective lepton mixing matrix U is given by

$$U = V^\dagger U_{13}\, {\rm e}^{i\frac{\pi}{2}} U_p \tag{12}$$

with

$$U_{13} = \begin{pmatrix} \cos\theta & 0 & {\rm e}^{-i\psi}\sin\theta \\ 0 & 1 & 0 \\ -{\rm e}^{i\psi}\sin\theta & 0 & \cos\theta \end{pmatrix} \tag{13}$$

and $U_p = {\rm diag}(1, {\rm e}^{i\beta_{21}/2}, {\rm e}^{i\beta_{31}/2})$. The parameters θ, ψ appearing in the unitary rotation matrix U_{13} can now be written as

$$\tan 2\theta = \frac{2\chi_1}{2\chi_1\chi_2\cos\gamma_2\cos\psi - \left[\chi_1^2\sin\gamma_1 + \chi_2^2\sin(2\gamma_2-\gamma_1) - \sin\gamma_1\right]\sin\psi}\,, \tag{14}$$

$$\tan\psi = \frac{-2\chi_1\chi_2\sin\gamma_2}{\cos\gamma_1 + \chi_2^2\cos(2\gamma_2-\gamma_1) + \chi_1^2\cos\gamma_1}\,, \tag{15}$$

where we have defined the parameters χ_1 and χ_2 as $\chi_1 = |f_1^\nu/f_3^\nu|$, $\chi_2 = |f_2^\nu/f_3^\nu|$, and the associated phase differences are defined as $(\phi_1 - \phi_3) = \gamma_1$, $(\phi_2 - \phi_3) = \gamma_2$. The real positive light-neutrino masses can be written as

$$m_1 = \frac{v^2}{M}\left|f_3^{\nu^2}\right|\sqrt{o_1^2 + o_2^2}\,, \tag{16}$$

$$m_2 = \frac{v^2}{M}\left|f_3^{\nu^2}\right|, \tag{17}$$

$$m_3 = \frac{v^2}{M}\left|f_3^{\nu^2}\right|\sqrt{n_1^2 + n_2^2}\,, \tag{18}$$

where o_1, o_2, n_1, and n_2 in our model take the form of

$$o_1 = \chi_1^2 \cos 2\gamma_1 + \chi_2^2 \cos 2\gamma_2 + 1 - 2A\chi_1 \cos\gamma_1 + 2B\chi_1 \sin\gamma_1\,, \tag{19}$$

$$o_2 = \chi_1^2 \sin 2\gamma_1 + \chi_2^2 \sin 2\gamma_2 - 2A\chi \sin\gamma_1 - 2B\chi_1 \cos\gamma_1\,, \tag{20}$$

$$n_1 = \chi_1^2 \cos 2\gamma_1 + \chi_2^2 \cos 2\gamma_2 + 1 + 2A\chi_1 \cos\gamma_1 - 2B\chi_1 \sin\gamma_1\,, \tag{21}$$

$$n_2 = \chi_1^2 \sin 2\gamma_1 + \chi_2^2 \sin 2\gamma_2 + 2A\chi \sin\gamma_1 + 2B\chi_1 \cos\gamma_1\,, \tag{22}$$

$$A = \frac{\sqrt{1 + \chi_2^2 \cos 2\gamma_2 + \sqrt{1 + \chi_2^4 + 2\chi_2^2 \cos 2\gamma_2}}}{\sqrt{2}}\,, \qquad B = \frac{\chi_2^2 \sin 2\gamma_2}{2A}\,. \tag{23}$$

The phases $\beta_{21(31)}$ involved in U_p are given by

$$\beta_{21} = -\tan^{-1}\frac{o_2}{o_1}\,, \qquad \beta_{31} = \tan^{-1}\frac{n_2}{n_1} - \tan^{-1}\frac{o_2}{o_1}\,. \tag{24}$$

Comparing the mixing matrix given in Eq. (12) with the standard form of the lepton mixing matrix ($U_{\rm PMNS}$) [27], we obtain the correlation between the neutrino mixing angles (and the Dirac CP phase) appearing in $U_{\rm PMNS}$ and the model parameters as [15]

$$|\, s_{13}\,|^2 = \frac{1 + \sin 2\theta \cos\psi}{3}\,, \qquad \tan\delta = \frac{\sin\theta \sin\psi}{\cos\theta + \sin\theta\cos\psi}\,, \tag{25}$$

$$s_{12}^2 = \frac{1}{3(1 - |\, s_{13}\,|^2)}\,, \quad \tan 2\theta_{23}\cos\delta = \frac{1 - 2\,|\, s_{13}\,|^2}{|\, s_{13}\,|\sqrt{2 - 3\,|\, s_{13}\,|^2}}\,, \tag{26}$$

and the two Majorana phases α_{21} and α_{31} are found to be $\alpha_{21} = \beta_{21}$, and $\alpha_{31} = \beta_{31}$.

3. Constraints neutrino masses and mixing

From Eqs. (14), (15), (25), and (26), it is clear that all three mixing angles $\theta_{13,12,23}$ and the Dirac CP phase (δ) appearing in $U_{\rm PMNS}$ can be determined by four model parameters χ_1, χ_2, γ_1, and γ_2. Let us now define a dimensionless quantity r as the ratio of solar (Δm^2_{31}) to atmospheric (Δm^2_{21}) mass squared difference

$$r = \frac{\Delta m^2_{21}}{\Delta m^2_{31}} = \frac{m^2_2 - m^2_1}{m^2_2 - m^2_1} = \frac{1 - o^2_1 - o^2_2}{n^2_1 + n^2_2 - o^2_1 - o^2_2}\,, \tag{27}$$

where we have used Eqs. (16)–(18). Furthermore, we can compute the absolute neutrino masses in Eqs. (16)–(18) by computing the pre-factor $|f^{\nu^2}_3| v^2/M$ via

$$\left|f^{\nu^2}_3\right| v^2/M = \sqrt{\frac{\Delta m^2_{21}}{(1 - o^2_1 - o^2_2)}}\,. \tag{28}$$

Therefore, just like the mixing angles, r, $m_{1,2,3}$, $\Delta m^2_{21,31}$, and the sum of the light-neutrinos masses $\sum_i m_i (\leq 0.11$ eV [28]) are also found to be a function of the same model parameters χ_1, χ_2, γ_1, and γ_2. Now, using the 3σ allowed range [29] for oscillation data and absolute mass [28], we can obtain the allowed ranges for these parameters as given in Fig. 1. Here, the blue dots represent the allowed regions in the χ_1–γ_1 plane, whereas the magenta dots represent allowed regions in the χ_2–γ_2 plane which stand for the antisymmetric contribution in Dirac Yukawa mentioned earlier. From Fig. 1, we find that when $0.584 \lesssim \chi_1 \lesssim 1.462$, γ_1 falls in a range of 0°–69° and 291°–360°, whereas when $0.470 \gtrsim \chi_2 \gtrsim 0.145$, γ_2 falls within 78°–161° and 206°–287°, respectively. In our analysis, the Dirac CP phase is found to be $33°(213°) \lesssim \delta \lesssim 80°(260°)$ and $100°(280°) \lesssim \delta \lesssim 147°(327°)$ which corresponds to the atmospheric mixing angle θ_{23} in the lower octant dictated by correlation given in Eq. (26). Similarly, in our analysis, the lower limit on the sum of the light-neutrino mass $\sum_i m_i$ is found to be 0.06 eV. The present analysis strictly predicts the normal hierarchy of light-neutrino mass

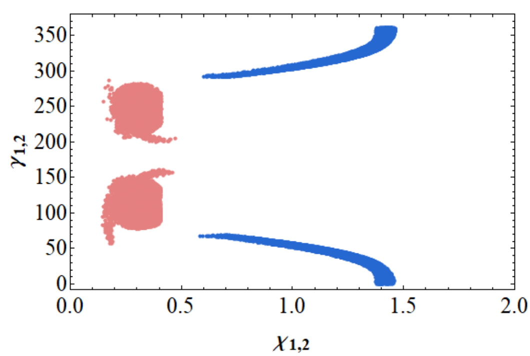

Fig. 1. Allowed parameter spaces of χ_1–γ_1 (blue dots) and χ_2–γ_2 (magenta dots) using 3σ ranges of neutrino oscillation parameters [29].

as a consequence of the adopted flavor symmetry. This is indeed a distinctive and important prediction. Now, from Eq. (24) and the relations $\alpha_{21} = \beta_{21}$, and $\alpha_{31} = \beta_{31}$, we can also obtain the allowed region for the Majorana phases as given in the left panel of Fig. 2. Interestingly, it is important to note that the Majorana phases cannot be constrained by the neutrino oscillation experiments. However, here we were able to constrain them using low-energy oscillation data. After computing the Majorana phases (as we have already evaluated $\theta_{12,13}$, $m_{1,2,3}$, and δ earlier), we are now in a position to make prediction for the effective mass parameter appearing in the neutrinoless double beta decay as $0.002 \lesssim m_{\beta\beta} \lesssim 0.021$ eV for the normal hierarchy of light-neutrino mass as given in the right panel of Fig. 2 (taken from [22]).

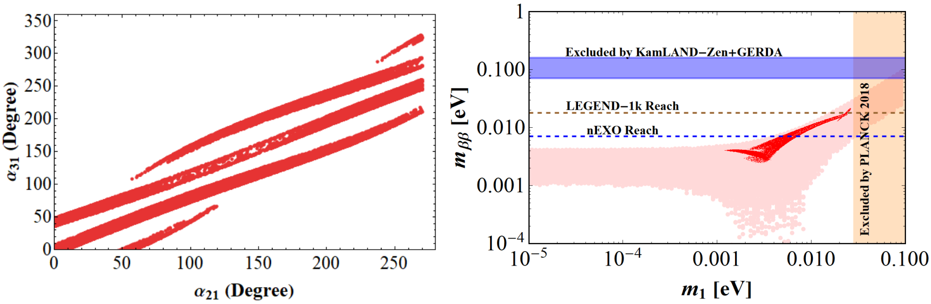

Fig. 2. Left panel: Prediction for the Majorana phases α_{21} and α_{31}. Right panel: Plot for $m_{\beta\beta}$ *versus* m_1 (for the normal hierarchy) with allowed ranges for χ_1, χ_2, γ_1, and γ_2 obtained from Fig. 1.

4. Leptogenesis

The out-of-equilibrium decay of heavy RHNs in the early Universe can produce lepton asymmetry [30]. The created lepton asymmetry can be converted to a baryon asymmetry via the sphaleron process. Taking the lepton flavor effect into consideration, the CP asymmetry parameter can be expressed as [31, 32]

$$\begin{aligned}\epsilon_i^\alpha &= \frac{1}{8\pi\mathcal{H}_{ii}}\sum_{j\neq i}\mathrm{Im}\left[\mathcal{H}_{ij}\left(\mathcal{Y}^{\nu\dagger}\right)_{i\alpha}\left(\mathcal{Y}^\nu\right)_{\alpha j}\right]\left[f(x_{ij}) + \frac{\sqrt{x_{ij}}(1-x_{ij})}{(1-x_{ij})^2 + \frac{\mathcal{H}_{jj}^2}{64\pi^2}}\right] \\ &\quad + \frac{1}{8\pi\mathcal{H}_{ii}}\sum_{j\neq i}\mathrm{Im}\left[\mathcal{H}_{ji}\left(\mathcal{Y}^{\nu\dagger}\right)_{i\alpha}\left(\mathcal{Y}^\nu\right)_{\alpha j}\right]\left[\frac{(1-x_{ij})}{(1-x_{ij})^2 + \frac{\mathcal{H}_{jj}^2}{64\pi^2}}\right], \qquad (29)\end{aligned}$$

where $\mathcal{Y}^\nu$ ($\equiv V^\dagger Y^\nu$) is the neutrino Yukawa matrix in the basis where the charged leptons are diagonal, and $\mathcal{H}$ and the loop factor $f(x_{ij})$ are given by

$$\mathcal{H} \; = \; \mathcal{Y}^{\nu\dagger}\mathcal{Y}^{\nu} = Y^{\nu\dagger}Y^{\nu}\,; \tag{30}$$

$$f(x_{ij}) \; = \; \sqrt{x_{ij}}\left[1-(1+x_{ij})\ln\left(\frac{1+x_{ij}}{x_ij}\right)\right]\,, \tag{31}$$

with $x_{ij} = \frac{M_j^2}{M_i^2}$, where M_i are the masses of the RHNs once the degeneracy is lifted. For strongly hierarchical RHNs, one can neglect $\frac{\mathcal{H}_{jj}^2}{64\pi^2}$ compared to $(1-x_{ij})^2$, while the entire expression of Eq. (29) can be used for quasi-degenerate regime even in resonance situation for which $(1-x_{ij})^2 \simeq \frac{\mathcal{H}_{jj}^2}{64\pi^2}$ [32, 33]. Here, we refer the readers to [19, 22] for a detailed study of the renormalization group effect to break the degeneracy of the RHNs to make leptogenesis viable in the present scenario. In Fig. 3 (taken from [22]), we plot variation of individual flavor components of CP asymmetry with respect to a model parameter χ_1 for three different scales: $M = 10^{13}$ GeV (top panel), $M = 10^{11}$ GeV (middle panels), and $M = 10^8$ GeV (bottom panels). In each case, depending on the temperature, flavor effects come into play and

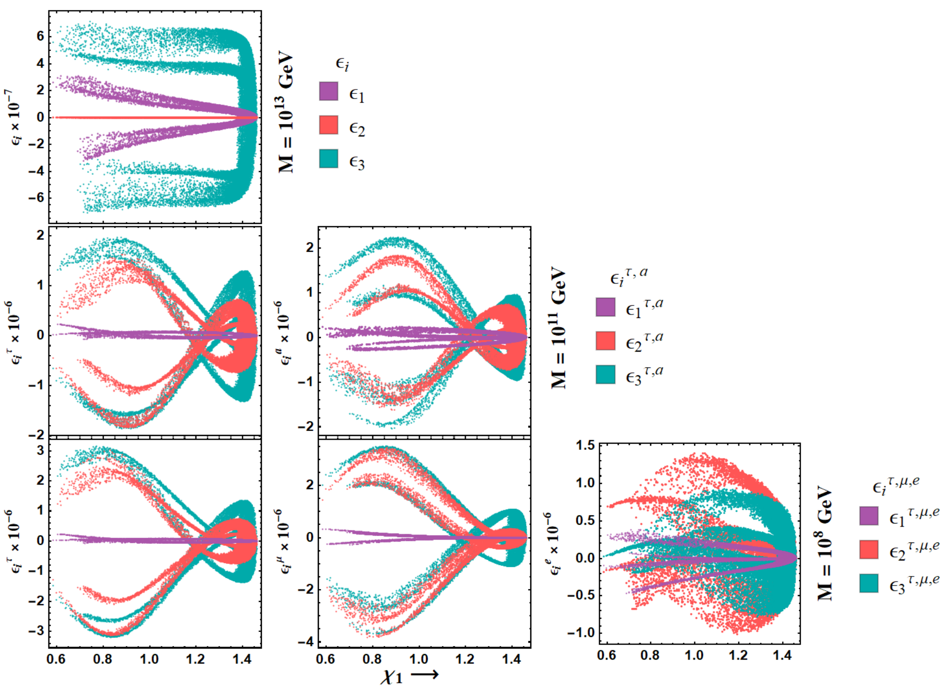

Fig. 3. Variation of CP asymmetries w.r.t. χ_1 for three different scales: $M = 10^{13}$ GeV (top panel), $M = 10^{11}$ GeV (middle panels), and $M = 10^8$ GeV (bottom panels).

generate CP asymmetry in the correct ballpark. Therefore, to compute the final baryon asymmetry, we need to solve the relevant Boltzmann equation taking the flavor effect into consideration. For a system of three RHNs, these equations can be written as [34–36]

$$
\begin{aligned}
sHz\frac{\mathrm{d}Y_{N_i}}{\mathrm{d}z} &= -\Bigg\{\left(\frac{Y_{N_i}}{Y^{\mathrm{eq}}_{N_i}}-1\right)\left(\gamma_{D_i}+2\gamma_{N^i_s}+4\gamma_{N^i_t}\right)\\
&\quad+\sum_{j\neq i}\left(\frac{Y_{N_i}}{Y^{\mathrm{eq}}_{N_i}}\frac{Y_{N_j}}{Y^{\mathrm{eq}}_{N_j}}-1\right)\left(\gamma^{(1)}_{N_iN_j}+\gamma^{(2)}_{N_iN_j}\right)\Bigg\}\,, \qquad (32)\\
sHz\frac{\mathrm{d}Y_{\Delta_\alpha}}{\mathrm{d}z} &= -\Bigg\{\sum_i\left(\frac{Y_{N_i}}{Y^{\mathrm{eq}}_{N_i}}-1\right)\epsilon^\alpha_i\gamma_{D_i}-\sum_\beta\Bigg[\sum_i\Bigg(\frac{1}{2}\left(C^\ell_{\alpha\beta}-C^H_\beta\right)\gamma^\alpha_{D_i}\\
&\quad+\left(C^\ell_{\alpha\beta}\frac{Y_{N_i}}{Y^{\mathrm{eq}}_{N_i}}-\frac{C^H_\beta}{2}\right)\gamma_{N^i_s}+\left(2C^\ell_{\alpha\beta}-\frac{C^H_\beta}{2}\left(1+\frac{Y_{N_i}}{Y^{\mathrm{eq}}_{N_i}}\right)\right)\gamma_{N^i_t}\Bigg)\\
&\quad+\sum_\gamma\Bigg(\left(C^\ell_{\alpha\beta}+C^\ell_{\gamma\beta}-2C^H_\beta\right)\left(\gamma^{(1)\alpha\gamma}_N+\gamma^{(2)\alpha\gamma}_N\right)\\
&\quad\sum_{i,j}\left(C^\ell_{\alpha\beta}-C^\ell_{\gamma\beta}\right)\gamma^{(1)\alpha\gamma}_{N_iN_j}\Bigg)\Bigg]\frac{Y_{\Delta_\beta}}{Y^{\mathrm{eq}}}\Bigg\}\,, \qquad (33)
\end{aligned}
$$

where $z = M_i/T$ and $\alpha = e, \mu, \tau$. Here, $Y_{\Delta_\alpha(N_i)} = n_{\Delta_\alpha(N_i)}/s$ denotes the density of $\Delta_\alpha = \frac{B}{3} - L_\alpha$ (relevant heavy neutrino) with respect to the entropy s, $C^{\ell,H}$ are the coefficient matrices which take care of the flavor effect at different temperatures, Y^{eq}s are the respective number densities while in thermal equilibrium. Here, the total decay rate density of N_i is given by

$$
\gamma_{D_i} = \sum_\alpha\left[\gamma\left(N_i\to\ell_\alpha+H\right)+\gamma\left(N_i\to\bar{\ell}_\alpha+\bar{H}\right)\right] = n^{\mathrm{eq}}_{N_i}\frac{K_1(z)}{K_2(z)}\mathit{\Gamma}_i\,, \qquad (34)
$$

where $\mathit{\Gamma}_i$ is the total decay rate of N_i at tree level and can be written as

$$
\mathit{\Gamma}_i = \sum_\alpha\left[\mathit{\Gamma}\left(N_i\to\ell_\alpha+H\right)+\mathit{\Gamma}\left(N_i\to\bar{\ell}_\alpha+\bar{H}\right)\right]\,, \qquad (35)
$$

and $\gamma_{N^i_s}$, $\gamma_{N^i_t}$, $\gamma^{(1)}_{N_iN_j}$, $\gamma^{(2)}_{N_iN_j}$ are the reaction rate densities for the scattering processes: $[N_i + \ell \leftrightarrow Q + \bar{U}\,]_s$, $[N_i + \bar{Q} \leftrightarrow \bar{\ell} + \bar{U}\,]_t + [N_i + U \leftrightarrow \bar{\ell} + \bar{Q}\,]_t$, $[N_i + N_j \leftrightarrow \ell + \bar{\ell}\,]$ and $[N_i + N_j \leftrightarrow H + \bar{H}\,]$, respectively [37, 38], $K_1(z)$ and

$K_2(z)$ are the modified Bessel functions. Finally, everything now in hand, for two unique benchmark points BP1(★) = $(1.37, 0.399,\ 21.53^\circ, 135.59^\circ)$ and BP2(▲) = $(0.978, 0.235,\ 301.81^\circ, 119.1^\circ)$, we plot the results in Fig. 4 (obtained from [22]). In each of the panels of Fig. 4, we estimate the evolution of the B–L asymmetry (denoted by red dotted line) as well as the final B asymmetry (denoted by magenta solid line) for the considered benchmark points BP1 (left panel), BP2 (right panel), respectively. Contributions coming from each flavor component are also shown in these plots. The horizontal cyan patch in both panels represents the observed value of the baryon asymmetry and for the considered benchmark one can satisfy this observation.

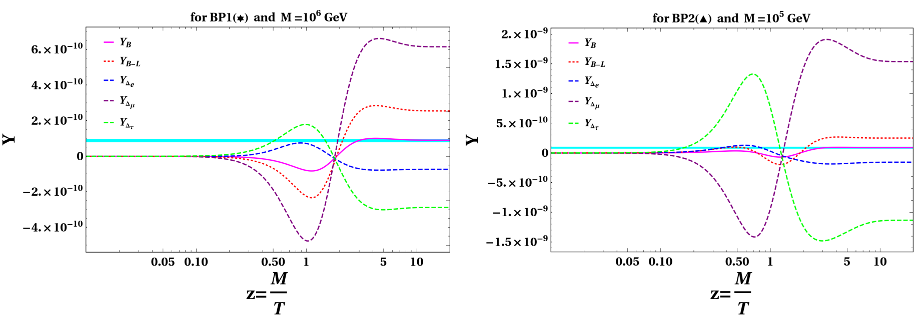

Fig. 4. Variation of Y_B, Y_{B-L}, Y_{Δ_e}, Y_{Δ_μ}, $Y_{\Delta\tau}$ (denoted by solid magenta, dotted red, dashed blue, dashed pink, and dashed green lines, respectively) presented as a function of $z = M/T$. Here, we have considered two benchmark points, namely BP1 (left panel) and BP2 (right panel) from Fig. 1.

5. Conclusions

In this work, we have studied a type-I seesaw model based on $A_4 \times Z_3 \times Z_2$ discrete symmetry to explain correct neutrino masses and mixing which also addresses the issues of leptogenesis. In terms of the flavon content, this model coincides with the original AF model (proposed to explain TBM mixing) with three flavon fields only. Thanks to the considered transformations of fields under the symmetry, here we are able to reproduce correct neutrino mixing data where an antisymmetric contribution in the Dirac Yukawa coupling plays a crucial role. After constraining the model parameters, we make important predictions for θ_{23} (found to be in the lower octant), δ (33°–80°, 100°–147°, 213°–260°, 280°–327°), the neutrino mass hierarchy (only the normal hierarchy is allowed), $\sum_i m_i (\gtrsim 0.06$ eV), and $m_{\beta\beta}$ (0.002–0.02 eV) as a consequence of the considered specific flavor symmetry. Another important finding in this particular set-up is that we can constrain the Majorana phases which are otherwise insensitive to oscillation experiments. Here, to start with, the RHNs are found to be exactly degenerate which forbids leptogenesis to take place. However, this can be elegantly solved by considering renormalization group effects which break this degeneracy, and leptogenesis becomes viable. Finally, we have obtained the observed baryon asymmetry of the universe by solving relevant Boltzmann equations taking lepton flavor effects into the picture.

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