PHYSICAL REVIEW D **89,** 096010 (2014)

Leptogenesis and CP violation in $SU(5)$ models with lepton flavor mixing originating from the right-handed sector

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(Received 13 January 2014; revised manuscript received 23 April 2014; published 27 May 2014)

We discuss neutrino masses and mixing in the context of seesaw type I models with three right-handed Majorana neutrinos and an approximately diagonal Dirac sector. This ansatz is motivated by the idea that the flavor structure in the right-handed Majorana masses is responsible for the large mixing angles, whereas the small mixing angle θ_{13} originates from the Dirac Yukawa couplings in analogy to the quark sector. To obtain $\theta_{13} \approx 0.15$ we study a possible $SU(5)$ grand unified theory realization with a $U(1) \times \mathbb{Z}_2' \times \mathbb{Z}_2'' \times \mathbb{Z}_2'''$ flavor symmetry and include a complex perturbation parameter in the Dirac mass matrix. The consequences for CP violating phases and effects on leptogenesis are investigated.

DOI: 10.1103/PhysRevD.89.096010 PACS numbers: 11.30.Hv, 12.10.-g, 14.60.St

I. INTRODUCTION

Neutrinos are special in several respects: First, they are much lighter than the charged leptons and quarks, and as they do not carry any unbroken quantum number they can be Majorana particles. These properties are exploited in the seesaw mechanism where the heavy Majorana masses of right-handed neutrinos drive the effective masses of left-handed neutrinos down to or below the eV scale. Moreover, the mismatch between neutrino and charged lepton mixing parametrized in the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix exhibits large or even maximal mixing angles, in stark contrast to the small Cabibbo-Kobayashi-Maskawa (CKM) mixing in the quark sector. After recent data from reactor experiments have revealed a nonvanishing neutrino mixing angle θ_{13} and thus rule out exact tribimaximal (TBM) mixing [1–3], many attempts have been made to explain the finite θ_{13}. While the idea of anarchy is becoming more attractive, the most popular approaches are still based on discrete symmetries, as summarized, e.g., in [4] or [5]. Models that intend to explain the leptonic mixing pattern can be constructed using large symmetry groups such as $\Delta(96)$ or D_{14}. At the expense of predictivity, these models contain several free parameters to account for all the physical observables. Alternatively, models can be based on smaller symmetry groups, e.g., A_4 or S_4, which yield specific patterns such as the TBM or golden ratio mixing and include perturbations in order to generate the necessary deviations from these structures.

In this paper we analyze the possibility that also the large lepton mixing arises from the right-handed Majorana sector. To this end we study a generic type I seesaw model with three heavy right-handed neutrinos,

$$m_\nu = m_D^T M_R^{-1} m_D, \tag{1}$$

where m_D denotes the Dirac mass matrix and M_R the matrix of the right-handed neutrinos. In a previous publication [6] the Dirac matrix has been assumed to be diagonal, and it has been shown that this can give rise to TBM. We adjust this pattern by perturbing the Dirac mass matrix with a complex parameter in order to accommodate a finite θ_{13}. The philosophy behind this ansatz is to have the small mixing originate from the Dirac sector, in analogy to the small CKM mixing in the quark sector, while the Majorana property of the right-handed neutrinos is responsible for the large mixing angles θ_{23} and θ_{12}. The structure of the Dirac sector is motivated by an $SU(5)$ Grand Unified Theory (GUT) embedding, which accounts not only for leptonic mixing but also for observables of the quark sector. Related works exist, where the Majorana mass terms of the neutrino sector account for the large leptonic mixing angles, e.g., [7,8].

The paper is organized as follows: We start by describing the outline of the model in Sec. II proposing a possible realization in an $SU(5)$ GUT. In Sec. III we briefly summarize the methods used to analyze CP violation in our model and then focus on leptogenesis in Sec. IV. Section V deals with the numerical analysis and with the implications for neutrinoless double-beta decay ($0\nu\beta\beta$). Conclusions are drawn in Sec. VI.

II. OUTLINE OF THE MODEL

If we consider m_D and the mass matrix of the charged leptons to be diagonal, all mixing in the neutrino sector originates from M_R. An explicit example can be constructed by using an $SU(5)$ GUT with $U(1) \times \mathbb{Z}_2' \times \mathbb{Z}_2'' \times \mathbb{Z}_2'''$ flavor symmetry in close analogy to the model published in [9]. All standard model particles, including the right-handed neutrinos, are accommodated in the $\mathbf{10} \oplus \mathbf{5} \oplus \mathbf{1}$ multiplets of $SU(5)$. Assigning appropriate

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quantum numbers under these symmetry groups yields the quark mixing matrix in the first approximation, where the off-diagonal elements are suppressed by a Froggatt-Nielsen (FN) mechanism [10]. In this framework the fermion masses are generated through couplings to additional scalar fields η_i with universal vacuum expectation values (VEVs) u. These VEVs are suppressed by a large messenger scale Λ such that $\frac{u}{\Lambda} \approx \lambda = 0.22$, consequently leading to suppression factors in the fermion mass term $\propto \lambda^n$, where n denotes the sum of the fermion field charges under the corresponding flavor symmetries. It has been shown, e.g., in [11,12] that the combination of an $SU(5)$ with a $U(1)_{\mathrm{FN}}$ symmetry can give rise to maximal mixing in the lepton sector due to specific $U(1)_{\mathrm{FN}}$ charge assignments.

The matrix structures generated by the FN mechanism must be consistent with the following approximate mass relations:

$$\begin{aligned} &m_u : m_c : m_t \approx \lambda^8 : \lambda^4 : 1, \qquad m_d : m_s : m_b \approx \lambda^4 : \lambda^2 : 1, \\ &m_e : m_\mu : m_\tau \approx \lambda^4 : \lambda^2 : 1. \end{aligned} \tag{2}$$

Following Ref. [9], assigning the following $U(1)$ charges to the $SU(5)$ multiplets, and (-1) for the flavon fields, leads to the desired structures of the Yukawa matrices in the quark sector,

$$\mathbf{10}\colon\ I:4 \quad II:2 \quad III:0, \tag{3}$$

$$\mathbf{5}\colon\ I:3 \quad II:3 \quad III:3, \tag{4}$$

where I, II, and III specify the family number. The powers of the suppression factors λ entering the Lagrangian depend on the $\mathbb{Z}_2$ charges ρ_1, ρ_2, and ρ_3 of the respective flavons η', η'', and η'''. The sum rules for cyclic groups $\mathbb{Z}_2$ are

$$\rho_i + \rho_j = (\rho_i + \rho_j) \mod 2, \tag{5}$$

where $\rho_{i,j}$ again are the $\mathbb{Z}_n$ charges of the interacting fermions $\psi_{i,j}$. This results in

$$\mathbf{10}_i \otimes \mathbf{10}_j\colon\ Y_u \propto \begin{pmatrix} \lambda^8 & \lambda^{6+\rho_1+\rho_2} & \lambda^{4+\rho_1+\rho_3} \\ \lambda^{6+\rho_1+\rho_2} & \lambda^4 & \lambda^{2+\rho_3+\rho_2} \\ \lambda^{4+\rho_1+\rho_3} & \lambda^{2+\rho_3+\rho_2} & 1 \end{pmatrix}, \tag{6}$$

$$\mathbf{10}_i \otimes \bar{\mathbf{5}}_j\colon\ Y_d \propto \begin{pmatrix} \lambda^7 & \lambda^{7+\rho_1+\rho_2} & \lambda^{7+\rho_1+\rho_3} \\ \lambda^{5+\rho_1+\rho_2} & \lambda^5 & \lambda^{5+\rho_3+\rho_2} \\ \lambda^{3+\rho_1+\rho_3} & \lambda^{3+\rho_3+\rho_2} & \lambda^3 \end{pmatrix}. \tag{7}$$

Since the Yukawa matrix of the charged leptons $Y_l \sim \bar{\mathbf{5}}_i \otimes \mathbf{10}_j \sim Y_d^T$ is hierarchical as well, large lepton mixing must be a consequence of the specific structure of the neutrino sector. To establish the structure of the latter, let us denote the $U(1)$ charges of the right-handed neutrinos $N^R_{1,2,3}$ in the $SU(5)$ singlet $\mathbf{1}$ as e_1, e_2, and e_3, respectively. According to Ref. [9] the right-handed neutrinos carry no $\mathbb{Z}_2$ charges as the seesaw scale is lower than the messenger scale Λ at which the flavons η_i receive their VEVs u. The relevant matrix structures in the neutrino sector arise from the following products:

$$\bar{\mathbf{5}}_i \otimes \mathbf{1}_j\colon\ Y^D \propto \begin{pmatrix} \lambda^{3+e_1} & \lambda^{3+e_2+\rho_1} & \lambda^{3+e_3+\rho_1} \\ \lambda^{3+e_1+\rho_2} & \lambda^{3+e_2} & \lambda^{3+e_3+\rho_2} \\ \lambda^{3+e_1+\rho_3} & \lambda^{3+e_2+\rho_3} & \lambda^{3+e_3} \end{pmatrix}, \tag{8}$$

$$\mathbf{1}_i \otimes \mathbf{1}_j\colon\ Y_R \propto \begin{pmatrix} \lambda^{2e_1} & \lambda^{e_1+e_2} & \lambda^{e_1+e_3} \\ \lambda^{e_2+e_1} & \lambda^{2e_2} & \lambda^{e_2+e_3} \\ \lambda^{e_3+e_1} & \lambda^{e_3+e_2} & \lambda^{2e_3} \end{pmatrix}. \tag{9}$$

By choosing

$$e_3 = 0; \qquad e_1 = e_2 = 1; \qquad \rho_1 = \rho_2 = \rho_3 = 1, \tag{10}$$

we can motivate a hierarchical structure of the Dirac matrix Y^D with two perturbation parameters assigned to the elements Y^D_{12} and Y^D_{23}:

$$Y^D \propto \lambda^3 \cdot \begin{pmatrix} \lambda & \lambda^2 & \lambda \\ \lambda^2 & \lambda & \lambda \\ \lambda^2 & \lambda^2 & 1 \end{pmatrix}, \qquad Y_R \propto \begin{pmatrix} \lambda^2 & \lambda^2 & \lambda \\ \lambda^2 & \lambda^2 & \lambda \\ \lambda & \lambda & 1 \end{pmatrix}. \tag{11}$$

Following our charge assignments from Eqs. (10) and (11) we choose a specific Dirac mass pattern for our further analysis,

$$m_D = \begin{pmatrix} \lambda & 0 & \epsilon \\ 0 & \lambda & \gamma \\ 0 & 0 & 1 \end{pmatrix} \cdot v, \tag{12}$$

with $\lambda \approx 0.22$ and $v = \frac{246}{\sqrt{2}}$ GeV, where ϵ generates a nonzero θ_{13} value and γ cancels the effects of ϵ on the large mixing angles θ_{12} and θ_{23}. In analogy to the quark sector we assume that all complex phases except for one can be absorbed into the interacting lepton fields. A single complex phase ϕ in $\epsilon = |\epsilon| e^{i\phi}$ remains, whereas $\gamma \in \mathbb{R}$.

The procedure to accommodate the experimentally determined mixing angles can be summarized in three steps. It has been demonstrated in Ref. [6] that a diagonal Dirac sector can give rise to TBM mixing. Therefore by adopting

$$m_D = \mathrm{diag}(\lambda, \lambda, 1) \cdot v \tag{13}$$

we first determine the right-handed mass matrix M_R to yield exact TBM mixing with θ_{12} and θ_{23} in the allowed 3σ ranges [13],

$$\begin{aligned} &\theta_{12} \in [0.543, 0.626], \qquad \theta_{23} \in [0.625, 0.956], \\ &\theta_{13} \in [0.125, 0.173]. \end{aligned} \tag{14}$$

We can derive an analytical expression for M_R^{-1} from the condition that m_ν must be diagonalizable with U_{TBM},

$$M_R^{-1} = (m_D^T)^{-1} U_{\mathrm{TBM}} m'_\nu U_{\mathrm{TBM}}^T m_D^{-1} \tag{15}$$

$$= \frac{1}{3v^2}\begin{pmatrix} \frac{2m_1+m_2}{\lambda^2} & \frac{-m_1+m_2}{\lambda^2} & \frac{-m_1+m_2}{\lambda} \\ \frac{-m_1+m_2}{\lambda^2} & \frac{m_1+2m_2+3m_3}{2\lambda^2} & \frac{m_1+2m_2-3m_3}{2\lambda} \\ \frac{-m_1+m_2}{\lambda} & \frac{m_1+2m_2-3m_3}{2\lambda} & \frac{m_1+2m_2+3m_3}{2} \end{pmatrix} \tag{16}$$

and

$$M_R = \frac{v^2}{3}\begin{pmatrix} \frac{\lambda^2(m_1+2m_2)}{m_1m_2} & \frac{\lambda^2(m_1-m_2)}{m_1m_2} & \frac{\lambda(m_1-m_2)}{m_1m_2} \\ \frac{\lambda^2(m_1-m_2)}{m_1m_2} & \frac{1}{2}\lambda^2\left(\frac{1}{m_1}+\frac{2}{m_2}+\frac{3}{m_3}\right) & \frac{1}{2}\lambda\left(\frac{1}{m_1}+\frac{2}{m_2}-\frac{3}{m_3}\right) \\ \frac{\lambda(m_1-m_2)}{m_1m_2} & \frac{1}{2}\lambda\left(\frac{1}{m_1}+\frac{2}{m_2}-\frac{3}{m_3}\right) & \frac{1}{2}\left(\frac{1}{m_1}+\frac{2}{m_2}+\frac{3}{m_3}\right) \end{pmatrix}, \tag{17}$$

where $m'_\nu = \mathrm{diag}(m_1, m_2, m_3)$ serves as an input parameter. The matrix M_R given in Eq. (17) is in good agreement with Y_R of Eq. (11). The initial values for the light neutrino masses m_i in m'_ν are selected according to the experimental bounds on Δm_{12}^2 and Δm_{23}^2 [14]

$$\Delta m_{12}^2 = 7.59\times 10^{-5}\ (\mathrm{eV})^2,$$
$$\Delta m_{23}^2 = \begin{cases} +2.46\times 10^{-3}\ (\mathrm{eV})^2\ (\mathrm{NH}) \\ -2.37\times 10^{-3}\ (\mathrm{eV})^2\ (\mathrm{IH}) \end{cases}. \tag{18}$$

The model exhibits different behavior in the various mass regimes, where the following cases are considered:

(1) Degenerate masses (deg): $m_1 \approx m_2 \approx m_3 > 0$,
(2) Normal hierarchy (NH): $m_3 \gg m_2 \approx m_1 \approx 0$,
(3) Inverse hierarchy (IH): $m_1 \approx m_2 \gg m_3 \approx 0$.

In the second step we include the perturbation parameters ϵ and γ in the Dirac mass matrix m_D according to Eq. (12) to produce $\theta_{13} \approx 0.15 = \theta_{13}^{\mathrm{Exp}}$, while M_R remains unchanged. This guarantees that θ_{13} is solely affected by Dirac couplings.

The parameters $|\epsilon|$, ϕ, and γ then are fitted to the current experimental 3σ bounds on the mixing angles θ_{ij} given in Eq. (14) for neutrino mass eigenstates $m_i \in [10^{-3}, 10^{-1}]$ eV. This mass region includes the NH and IH scenarios as well as the degenerate mass regime being consistent with cosmological bounds on the neutrino masses [15] and a successful leptogenesis scenario [16]. The bounds on the mixing angles constrain the parameter space of ϵ to small regions, enabling predictions on the observable CP phases. Assuming that leptogenesis successfully generates the correct baryon asymmetry of the Universe, additional constraints arise that confine these intervals even further.

III. CP VIOLATION

The PMNS matrix depends on four parameters: three mixing angles θ_{ij} and one CP violating phase δ. In the case of three additional right-handed neutrinos the mixing matrix includes two supplementary Majorana phases, leading to the conventional parametrization of the PMNS matrix

$$U_{\mathrm{PMNS}} = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta} & c_{23}c_{13} \end{pmatrix}\cdot P, \tag{19}$$

where $P = \mathrm{diag}(e^{i\alpha}, e^{i\beta}, 1)$ contains the Majorana phases α and β, and s_{ij} and c_{ij} are abbreviations for $\sin\theta_{ij}$ and $\cos\theta_{ij}$, respectively.

In our model the complex perturbation parameter ϵ provides the source of CP violation in the neutrino sector and leads to CP violating phases in U. The mixing matrix can be extracted from $m'_\nu = U^T m_\nu U$ using a Takagi decomposition, which is applicable for complex symmetric matrices [17]. The real diagonal matrix m'_ν contains the non-negative square roots of the eigenvalues of $m_\nu m_\nu^\dagger$. However, the mixing matrix we receive from our numerical analysis does not resemble the usual convention of the PMNS matrix given in Eq. (19). We can extract the mixing angles θ_{ij} and CP phases δ, α, and β using rephasing invariants and properties of the unitarity triangles of U, previously explored in [18,19]. By defining “Majorana-type” phases ζ_i and ξ_i, $i \in (1,2,3)$,

$$\zeta_1 \equiv \mathrm{Arg}(U_{e1}U_{e2}^*), \qquad \zeta_2 \equiv \mathrm{Arg}(U_{\mu1}U_{\mu2}^*), \qquad \zeta_3 \equiv \mathrm{Arg}(U_{\tau1}U_{\tau2}^*), \tag{20}$$

$$\xi_1 \equiv \mathrm{Arg}(U_{e1}U_{e3}^*), \qquad \xi_2 \equiv \mathrm{Arg}(U_{\mu1}U_{\mu3}^*), \qquad \xi_3 \equiv \mathrm{Arg}(U_{\tau1}U_{\tau3}^*), \tag{21}$$

we can express the mixing angles and δ according to [20] as a function of the following phases:

$$\tan^2\theta_{12} = \frac{|\sin(\xi_1-\xi_2)||\sin(-\zeta_2+\xi_2+\zeta_3-\xi_3)||\sin(\xi_1-\xi_3)|}{|\sin(-\zeta_1+\xi_1+\zeta_2-\xi_2)||\sin(\xi_2-\xi_3)||\sin(-\zeta_1+\xi_1+\zeta_3-\xi_3)|}, \tag{22}$$

$$\tan^2\theta_{23} = \frac{|\sin(\xi_1-\xi_3)||\sin(-\zeta_1+\xi_1+\zeta_3-\xi_3)||\sin(\zeta_1-\zeta_2)|}{|\sin(-\zeta_1+\xi_1+\zeta_2-\xi_2)||\sin(\xi_1-\xi_2)||\sin(\zeta_1-\zeta_3)|}, \tag{23}$$

$$\tan^2\theta_{13} = \frac{|\sin(\xi_2-\xi_3)||\sin(\zeta_1-\zeta_3)||\sin(\zeta_1-\zeta_2)|}{|\sin(\xi_1-\xi_3)||\sin(\xi_1-\xi_2)||\sin(\zeta_2-\zeta_3)|}\cdot\sin^2\theta_{12}, \tag{24}$$

$$\frac{1}{8}|\sin\delta| = \frac{|\cos\theta_{12}\cos\theta_{13}\sin\theta_{13}|^2|\sin(\xi_1-\xi_3)||\sin(\xi_1-\xi_2)|}{|\sin(2\theta_{12})\sin(2\theta_{13})\sin(2\theta_{23})\cos\theta_{13}||\sin(\xi_2-\xi_3)|}. \tag{25}$$

These formulas are valid regardless of the parametrization of the matrix U. To extract the Majorana phases α and β we recall that they rotate the Majorana unitarity triangles in the complex plane. Hence, one can conclude that expressions for the phases α and β result from (with $U_{\mathrm{PMNS},ij} \coloneqq U_{Pij}$)

$$\mathrm{Arg}\left(\frac{U_{P11}U^*_{P13}}{U_{11}U^*_{13}}\right) = \alpha, \qquad \mathrm{Arg}\left(\frac{U_{P12}U^*_{P13}}{U_{12}U^*_{13}}\right) = \beta,$$
$$\mathrm{Arg}\left(\frac{U_{P11}U^*_{P12}}{U_{11}U^*_{12}}\right) = \alpha-\beta. \tag{26}$$

IV. LEPTOGENESIS

Leptogenesis [21] explains the present matter asymmetry by assuming that a lepton asymmetry in the early universe is converted into a baryon asymmetry through $B+L$ violating sphaleron processes. Leptogenesis is closely connected to the seesaw mechanism as it relies on the existence of heavy right-handed Majorana neutrinos N_i decaying into leptons l_α and scalar fields ϕ via $N_i \to l_\alpha + \phi$. Observations of the cosmic microwave background allow for an estimate of the baryon asymmetry [22]

$$Y^{\mathrm{CMB}}_{\Delta\mathrm{B}} = \frac{n_B - n_{\bar{B}}}{s} = (8.79 \pm 0.44)\times 10^{-11}, \tag{27}$$

where s is the entropy density of the Universe and $n_{B,\bar{B}}$ denote the abundances of baryons and antibaryons, respectively. The baryon asymmetry can be calculated in terms of the CP asymmetry $\sigma_{i\alpha}$ and the efficiency factor κ_i [23],

$$Y_{\Delta\mathrm{B}} = \sigma_{i\alpha}\kappa_i \times 10^{-3}. \tag{28}$$

The efficiency factor κ_i measures the effect of washout processes depending on the "washout regime." For a general system it is given by the solution of the Boltzmann equations for leptogenesis, which characterize the competition of production and washout of the N_i's. In this paper the following approximations are used, as described in Ref. [23]:

$$\tilde{m} < m_* \cap m_i < m_*:\ \kappa_i \approx \frac{m_i\tilde{m}}{m_*^2} \quad \text{weak washout}, \tag{29}$$

$$\tilde{m} > m_* \cap m_i < m_*:\ \kappa_i \approx \frac{m_i}{m_*} \quad \text{intermediate washout}, \tag{30}$$

$$\tilde{m} > m_* \cap m_i > m_*:\ \kappa_i \approx \frac{m_*}{m_i} \quad \text{strong washout}, \tag{31}$$

with $m_* \approx 1.1\times 10^{-3}$ eV and the sum of the light neutrino masses $\tilde{m} = \sum_i m_i$.

The CP asymmetry $\sigma_{i\alpha}$ generated by a heavy right-handed neutrino N_i that decays into a lepton with flavor α reads explicitly [23]

$$\sigma_{i\alpha} \equiv \frac{1}{8\pi}\frac{1}{(Y^{D\dagger}Y^D)_{ii}}\sum_{j\neq i}\left\{\mathrm{Im}[(Y^{D\dagger}Y^D)_{ji}Y_{\alpha i}Y^{D*}_{\alpha j}]g\left(\frac{M_j^2}{M_i^2}\right)\right.$$
$$\left. + \mathrm{Im}[(Y^{D\dagger}Y^D)_{ij}Y^D_{\alpha i}Y^{D*}_{\alpha j}]\left(\frac{M_i^2}{M_i^2-M_j^2}\right)\right\}. \tag{32}$$

The Yukawa couplings matrix Y^D is related to the Dirac matrix by $Y^D = \frac{m_D}{v}$. Because of the Majorana nature of the N_i, contributions to the CP asymmetry arise only from higher-order interferences of their decays. The loop corrections are included in the function $g(x)$,

$$g(x) = \sqrt{x}\left[\frac{1}{1-x} + 1 - (1+x)\ln\left(1+\frac{1}{x}\right)\right]. \tag{33}$$

In a model with three additional right-handed neutrinos, the neutralizing effect of the $N_{2,3}$ on the asymmetry generated in the decay of N_1 can be neglected if $M_{2,3} > T_{\mathrm{reh}}$, the reheating temperature, and $M_1 \ll M_{2,3}$. In our approximation the second term in Eq. (32) vanishes due to zeros in the pattern of Y^D.

Model independent limits on neutrino masses can be inferred from the upper bound of the CP asymmetry $\sigma_{i\alpha}$. As stated in Sec. II, in the single lepton flavor approximation with hierarchical right-handed neutrino masses M_i a successful leptogenesis mechanism is confined to the mass region 10^{-3} eV $< m_i < 0.1$ eV [16]. We will therefore focus our attention on this interval.

Because of the specific structure of m_D given in Eq. (12), we obtain a small number of contributions to the CP asymmetry. With

$$Y^D = \begin{pmatrix} \lambda & 0 & |\epsilon|e^{i\phi} \\ 0 & \lambda & \gamma \\ 0 & 0 & 1 \end{pmatrix}, \quad Y^{D*} = \begin{pmatrix} \lambda & 0 & |\epsilon|e^{-i\phi} \\ 0 & \lambda & \gamma \\ 0 & 0 & 1 \end{pmatrix}, \quad \text{and}$$

$$Y^\dagger Y = \begin{pmatrix} \lambda^2 & 0 & |\epsilon|e^{i\phi}\lambda \\ 0 & \lambda^2 & \gamma\lambda \\ |\epsilon|e^{-i\phi}\lambda & \gamma\lambda & 1 \end{pmatrix}, \tag{34}$$

we receive nonvanishing contributions to $\sigma_{i\alpha}$ only if $\alpha = 1$,

$$\sigma_1 = \sum_i \sigma_{i1} = \frac{|\epsilon|^2 \sin(2\phi)}{8\pi} \left(\lambda^2 \cdot g\left(\frac{M_1^2}{M_3^2}\right) - g\left(\frac{M_3^2}{M_1^2}\right)\right), \tag{35}$$

resulting in a $|\epsilon|^2 \sin(2\phi)$ dependence of $Y_{\Delta \mathrm{B}}$. Note that the parameter γ does not affect the generated *CP* asymmetry.

V. NUMERICAL ANALYSIS

For the numerical analysis we use the following ordering schemes:

$$\text{NH}: \; m_1 = m_0, \qquad m_2 = \sqrt{m_0^2 + \Delta m_{12}^2}, \qquad m_3 = \sqrt{m_0^2 + \Delta m_{12}^2 + \Delta m_{23}^2}, \tag{36}$$

$$\text{IH}: \; m_3 = m_0, \qquad m_2 = \sqrt{m_0^2 + \Delta m_{12}^2 + \Delta m_{23}^2}, \qquad m_1 = \sqrt{m_0^2 + \Delta m_{23}^2}, \tag{37}$$

where $m_0 \in [10^{-3}, 10^{-1}]$ eV. In the regions of larger masses, where $m_i \gg \sqrt{\Delta m_{12}^2}$, $\sqrt{\Delta m_{23}^2}$ the mass eigenstates are degenerate. To explain the small neutrino masses through heavy right-handed neutrinos, which are compatible with the leptogenesis mechanism described in Sec. IV, the Dirac masses have to be at the GeV scale. As explained in Sec. I, we assume TBM mixing with a diagonal Dirac sector in order to determine M_R from Eq. (15).

In the cases of $m_0 = 10^{-3}$ eV, where the NH and IH scenarios are relevant, and $m_0 = 0.1$ eV (degenerate masses) the right-handed mass eigenstates are given by

$$\text{NH}: \; M_{\mathrm{R},1} = 4.934 \times 10^{13}\ \text{GeV}, \quad M_{\mathrm{R},2} = 3.942 \times 10^{14}\ \text{GeV}, \quad M_{\mathrm{R},3} = 7.311 \times 10^{15}\ \text{GeV}, \tag{38}$$

$$\text{IH}: \; M_{\mathrm{R},1} = 2.890 \times 10^{13}\ \text{GeV}, \quad M_{\mathrm{R},2} = 5.452 \times 10^{13}\ \text{GeV}, \quad M_{\mathrm{R},3} = 1.595 \times 10^{16}\ \text{GeV}, \tag{39}$$

$$\text{deg}: \; M_{\mathrm{R},1} = 1.365 \times 10^{13}\ \text{GeV}, \quad M_{\mathrm{R},2} = 1.447 \times 10^{13}\ \text{GeV}, \quad M_{\mathrm{R},3} = 2.830 \times 10^{14}\ \text{GeV}, \tag{40}$$

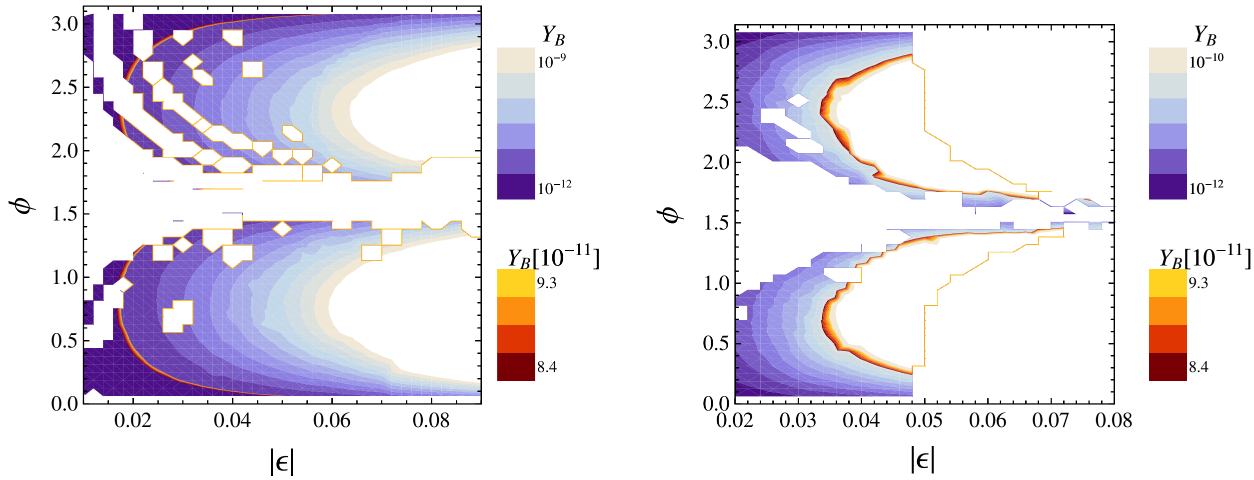
revealing a hierarchy among the the heavy right-handed neutrino masses, fulfilling the conditions for our approximations in leptogenesis.

The parameter spaces of the perturbation parameters $|\epsilon|$, ϕ, and γ are scanned for combinations that reproduce the neutrino mixing angles θ_{ij} within the current 3σ bounds; see Eq. (14). Assuming that leptogenesis successfully generates the correct baryon asymmetry of the Universe, given in Eq. (27), we can confine these regions even further to make precise predictions on physical observables.

The comparison of low-scale experimental data on neutrino parameters with calculations carried out at the GUT scale requires taking into account the renormalization group (RG) evolution of the leptonic mixing parameters. The RG running effects on the mixing parameters have been considered in various publications [24–26] and can be significant especially in the case of a degenerate mass spectrum. The leptonic mixing angles are expected to run faster than their quark equivalents for they are larger and the neutrino mass differences are particularly tiny. The generic enhancement for the evolution of the angles can be estimated analytically, which has been done, for instance, in [25]. For our numerical analysis of the degenerate mass regime we calculate the light neutrino masses at the seesaw scale using Eq. (1) and run the matrix m_ν down to the electroweak scale before determining the mixing parameters according to Sec. III. This allows for a better comparison with experimental data and a more robust analysis. The effects of RG running are implemented using REAP 1.8.4, generously provided by [25].

TABLE I. Parameter spaces of the fit parameters $|\epsilon|$, ϕ, and γ for successful leptogenesis and the *CP* phases δ, α, and β corresponding to these intervals.

	m_0	$\|\epsilon\|$	ϕ	γ
NH	[0.018, 0.050]	[0.018, 0.060]	$[0.063, 0.565] \cup [2.639, 3.079]$	[0.02, 0.20]
IH	[0.002, 0.044]	[0.034, 0.076]	$[0.251, 1.445] \cup [1.696, 2.890]$	$[-0.38, 0.00]$
		δ	α	β
NH		[0.016, 0.083]	$[0.039, 0.192] \cup [3.118, 3.133]$	[0.003, 0.157]
IH		[0.205, 1.293]	[0.448, 3.112]	$[0.009, 0.067] \cup [0.402, 2.404]$



(a) (b)

FIG. 1 (color online). (a) Parameter space of $|\epsilon|$ and ϕ compatible with the 3σ ranges of the θ_{ij} in the case of NH with $m_0 \in [0.00, 0.05]$ eV. The white areas did not yield models compatible with the experimental bounds. The (red) line overlying the contour plot marks parameter values that successfully generate the baryon asymmetry $Y_{\Delta B} = (8.79 \pm 0.44) \times 10^{-11}$ in leptogenesis. (b) Parameter space in the IH case.

Regarding leptogenesis, Eq. (31) is used to compute the efficiency factor κ_i since $m_0 > 10^{-3} > m^*$. The resulting parameter spaces complying with Eq. (27) are listed in Table I, where δ, α, β denote the Dirac and Majorana *CP* phases, respectively, corresponding to the parameter regions of $|\epsilon|$, ϕ, and γ. The allowed combinations of $|\epsilon|$ and ϕ as well as the generated baryon asymmetry are depicted in Fig. 1. The blue areas denote parameter values that account for neutrino mixing, while the red color indicates combinations that also lead to successful leptogenesis.

In a previous publication [6] it was found that models with a diagonal Dirac sector and TBM mixing favor very light neutrino masses. Although in principle all considered hierarchies are compatible with the experimental bounds on the neutrino mixing angles in our model, the statistics summarized in the following support the previous statement:

$$\text{NH}: \; m_0 \in [0.00, 0.05]\ \text{eV}, \quad n = 46250, \quad n_L = 19, \tag{41}$$

$$\text{IH}: \; m_0 \in [0.00, 0.05]\ \text{eV}, \quad n = 3768, \quad n_L = 229, \tag{42}$$

$$\text{deg}: \; m_0 \in [0.05, 0.10]\ \text{eV}, \quad n = 189, \quad n_L = 0, \tag{43}$$

where n and n_{L} count the number of events with and without leptogenesis, respectively, for each considered mass regime. The neutrino mixing is best accounted for by the NH scenario, whereas IH is favored if leptogenesis successfully explains the baryon asymmetry of the Universe. The RG analysis reveals that the model cannot accommodate the leptonic mixing angles for $m_0 > 0.056$ eV, practically ruling out degenerate neutrino masses as a suitable scenario.

Because of the number of combinations accommodating correct neutrino mixing as well as the baryon asymmetry of the early universe, predictions are the most reliable in the IH scenario. The predictions on the *CP* phases listed in Table I can be used for further studies; e.g., the Majorana phases α and β affect the observables measured in $0\nu\beta\beta$ experiments. The amplitude of these processes is proportional to the effective Majorana mass

$$m_{\beta\beta} = \sum_i U_{ei}^2 m_i. \tag{44}$$

Cancellation can occur in NH depending on the size of the Majorana phases; however, in this particular case we predict small low energy *CP* violation as can be seen from Table I. According to our results above we give an estimate of the effective Majorana mass in the considered hierarchies,

$$m_{\beta\beta}^{\mathrm{NH}} \in [0.048, 0.063]\ \text{eV}, \qquad m_{\beta\beta}^{\mathrm{IH}} \in [0.026, 0.050]\ \text{eV}. \tag{45}$$

Note again that due to Eqs. (41)–(43) the predictions are most reliable in the IH scenario. The bounds on $m_{\beta\beta}$ are well below the current upper limit $\langle m_{\beta\beta}\rangle \lesssim 0.19$–$0.45$ eV given by the EXO-200 experiment at 90% confidence level [27], but partly accessible in next-generation $0\nu\beta\beta$ decay experiments.

VI. CONCLUSIONS

$SU(5)$ inspired seesaw type I models with an almost diagonal Dirac sector and three heavy right-handed neutrinos successfully accommodate $\theta_{13}^{\mathrm{Exp}}$. The large neutrino mixing angles are ascribed to the structure of the right-handed sector, while the small mixing angle is generated by a complex perturbation parameter ϵ and a real parameter γ assigned to the off-diagonal elements of m_D. The structure of the perturbed Dirac mass matrix is obtained from an embedding in an $SU(5)\times U(1)\times \mathbb{Z}_2'\times \mathbb{Z}_2''\times \mathbb{Z}_2'''$ symmetry, where the hierarchy among the fermion families is generated by a Froggatt-Nielsen mechanism. The parameter ϵ also provides a possible source of CP violation in the leptonic sector, enabling predictions on the Dirac CP phase δ and the Majorana phases α and β.

The numerical analysis shows that the neutrino mixing parameters can in principle be reproduced in all considered hierarchies for masses up to $m_0 \approx 0.056$ eV. However, the NH scenario is strongly favored over IH and the degenerate mass regime accommodating the correct neutrino mixing angles θ_{ij}. Further constraints can arise from the requirement that leptogenesis generates the baryon asymmetry of the Universe, which eventually results in a preference of IH and rules out degenerate neutrino masses entirely. Regarding the phases δ, α, and β the CP violation in the IH case can be large, whereas for NH we predict small CP violation in all phases. Since the degenerate mass regime is excluded, the resulting effective Majorana mass $m_{\beta\beta}$ is partly accessible by next-generation $0\nu\beta\beta$ experiments in IH scenarios.

The leptogenesis scenario discussed in this paper is in good agreement with experimental and cosmological bounds, although a more realistic scenario with higher precision may improve the final predictions.

ACKNOWLEDGMENTS

H. P. was supported by DFG Grant No. PA 803/6-1. We thank Gautam Bhattacharyya for many helpful comments both on the original idea as well as on the actual work.

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