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Research Article

Cosmic Baryon Asymmetry in Different Neutrino Mass Models with Mixing Angles

Ng. K. Francis

Department of Physics, Tezpur University, Tezpur, Assam 784028, India

Correspondence should be addressed to Ng. K. Francis; ngkf2010@gmail.com

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We investigate the comparative studies of cosmological baryon asymmetry in different neutrino mass models with and without θ_{13} by considering the three-diagonal form of Dirac neutrino mass matrices and the three aspects of leptogenesis, unflavoured, flavoured, and nonthermal. We found that the estimations of any models with θ_{13} are consistent in all the three stages of calculations of leptogenesis and the results are better than the predictions of any models without θ_{13} which are consistent in a piecemeal manner with the observational data in all the three stages of leptogenesis calculations. For the normal hierarchy of Type-IA with charged lepton matrix, model with and without θ_{13} predicts inflaton mass required to produce the observed baryon asymmetry to be $M_\phi \sim 2.2 \times 10^{11}$ GeV and $M_\phi \sim 3.6 \times 10^{10}$ GeV, and the corresponding reheating temperatures are $T_R \sim 4.86 \times 10^6$ GeV and $T_R \sim 4.50 \times 10^6$ GeV respectively. These predictions are not in conflict with the gravitino problem which required the reheating temperature to be below 10^7 GeV. And these values apply to the recent discovery of Higgs boson of mass ~125 GeV. One can also have the right order of relic dark matter abundance only if the reheating temperature is bounded to below 10^7 GeV.

1. Introduction

Recent measurement of a moderately large value of the third mixing angle θ_{13} by reactor neutrino oscillation experiments around the world particularly by Daya Bay ($\sin^2\theta_{13} = 0.089 \pm 0.010(\text{stat}) \pm 0.005(\text{syst})$) [1] and RENO ($\sin^2\theta_{13} = 0.113 \pm 0.013(\text{stat}) \pm 0.019(\text{syst})$) [2] signifies an important breakthrough in establishing the standard three-flavour oscillation picture of neutrinos. Thereby, we will address the issues of the recent indication of nonmaximal 2-3 mixing by MINOS accelerator experiment [3] leading to determining the correct octant of θ_{23} and neutrino mass hierarchy. Furthermore, now, this has opened the door to study leptonic CP violation in a convincing manner, which in turn has profound implications for our understanding of the matter-antimatter asymmetry of the universe. In fact, ascertaining the origin of the cosmological baryon asymmetry, $\eta_B = (6.5^{+0.4}_{-0.5}) \times 10^{-10}$ [4], is one of the burning open issues in both particle physics and cosmology. The asymmetry must have been generated during the evolution of the universe. However, it is possible to dynamically generate such asymmetry if three conditions, (i) the existence of baryon number violating interactions, (ii) C and CP violations, and (iii) the deviation from thermal equilibrium, are satisfied [5]. There are different mechanisms of baryogenesis, but leptogenesis [6] is attractive because of its simplicity and the connection to neutrino physics. Establishing a connection between the low-energy neutrino mixing parameters and high-energy leptogenesis parameters has received much attention in recent years in [6–9]. In leptogenesis, the first condition is satisfied by the Majorana nature of heavy neutrinos and the sphaleron effect in the standard model (SM) at the high temperature [9], while the second condition is provided by their CP-violating decay. The deviation from thermal equilibrium is provided by the expansion of the universe. Needless to say the departures from thermal equilibrium have been very important without it; the past history of the universe would be irrelevant, as the present state would be merely that of a system at 2.75 K, very uninteresting indeed [10]. One of the keys to understanding the thermal history of the universe is the estimation of cosmological baryon asymmetry from different neutrino mass models with the inclusion of the latest nonzero θ_{13}.

Broadly the leptogenesis can be grouped into two groups: thermal with and without flavour effects and nonthermal leptogenesis. The simplest scenario, namely, the standard thermal leptogenesis, requires nothing but the thermal excitation of heavy Majorana neutrinos which generate tiny neutrino masses via the seesaw mechanism [11–13] and provides several implications for the light neutrino mass spectrum [14, 15]. And with heavy hierarchical right-handed neutrino spectrum, the CP asymmetry and the mass of the lightest right-handed Majorana neutrino are correlated. In order to have the correct order of light neutrino mass-squared differences, there is a lower bound on the mass of the right-handed neutrino, $M_N \geq 10^9$ GeV [16–19], which in turn put constraints on reheating temperature after inflation to be $T_R \geq 10^9$ GeV. This will lead to an excessive gravitino production and conflicts with the observed data. In the postinflation era, these gravitinos are produced in a thermal bath due to annihilation or scattering processes of different standard particles. The relic abundance of gravitino is proportional to the reheating temperature of the thermal bath. One can have the right order of relic dark matter abundance only if the reheating temperature is bounded to below 10^7 GeV [8, 20–24]. On the other hand, big-bang nucleosynthesis in SUSY theories also sets a severe constraint on the gravitino mass and the reheating temperature leading to the upper bound $T_R \geq 10^7$ GeV [25–29]. While thermal leptogenesis in SUSY SO(10) with high seesaw scale easily satisfies the lower bound, the tension with the gravitino constraint is manifest.

According to Fukuyama et al. [30, 31], the nonthermal leptogenesis scenario in the framework of a minimal supersymmetric SO(10) model with Type-I seesaw shows that the predicted inflaton mass needed to produce the observed baryon asymmetry of the universe is found to be $M_\phi \sim 5 \times 10^{11}$ GeV for the reheating temperature $T_R = 10^6$ GeV and weak scale gravitino mass $m_{3/2} \sim 100$ GeV without causing the gravitino problem. It also claims that even if these values are relaxed by one order of magnitude ($m_{3/2} \leq 10$ TeV, $T_R = 10^7$ GeV), the result is still valid. In [32, 33] using the Closed-Time-Path approach, they performed a systematic leading order calculation of the relaxation rate of flavour correlations of left-handed standard model leptons; and for flavoured leptogenesis in the early universe they found the reheating temperature to be $T_R = 10^7$ GeV to 10^{13} GeV. These values apply to the standard model with a Higgs-boson mass of 125 GeV [34]. The recent discovery of a standard model (SM) like Higgs boson provides further support for leptogenesis mechanism, where the asymmetry is generated by out-of-equilibrium decays of our conjecture heavy sterile right-handed neutrinos into a Higgs boson and a lepton. In [35] split neutrinos were introduced where there is one Dirac neutrino and two Majorana neutrinos with a slight departure from tribimaximal mixing (TBM), which explains the reactor angle $\sim \theta_{13}$, and tied intimately to the lepton asymmetry and can explain inflation, dark matter, neutrino masses, and the baryon asymmetry, which can be further constrained by the searches of SUSY particles at the LHC, the right-handed sneutrino, essentially the inflaton component as a dark matter candidate, and from the $0\nu\beta\beta$ experiments. In [36] too a deviation from TBM case was studied with model-independent discussion and the existing link between low- and high-energy parameters that connect to the parameters governing leptogenesis was analysed. However, in [37] exact TBM, $\tan^2\theta_{12} = 0.50$, was considered with charged lepton and up-quark type and set θ_{13} to be zero; eventually their results differ from ours. We slightly modify the neutrino models in [37]; consequently the inputs parameters are different for zero θ_{13} but for nonzero θ_{13} our formalism is entirely different than the one done in [37]; besides we consider $\tan^2\theta_{12} = 0.45$ for detail analysis. Our work in this paper is consistent with the values given in [30–35].

Now, the theoretical framework supporting leptogenesis from low-energy phases has some other realistic testable predictions in view of nonzero θ_{13}. So the present paper is a modest attempt to compare the predictions of leptogenesis from low-energy CP-violating phases in different neutrino mass matrices with and without θ_{13}. The current investigation is twofold. The first part deals with zero reactor mixing angle in different neutrino mass models within μ-τ symmetry [38–49], while in the second part we construct m_{LL} matrix from fitting of U_{PMNS} incorporating the nonzero third reactor angle (θ_{13}) along with the observed data and subsequently predict the baryon asymmetry of the universe (BAU). We must also mention that there are several works analysing the link between leptogenesis and low-energy data in more general scenarios. However, we have not come across in the literature where all the three categories of leptogenesis, that is, the thermal leptogenesis with or without flavour effects and nonthermal leptogenesis, are studied in a single paper. Take, for instance, some of the major players working on leptogenesis. Professor Wilfried Buchmuller works are mostly confined to standard unflavoured thermal leptogenesis by solving Boltzmann's equation whereas Professor Steven Blanchet and Professor P. Di. Bari generally worked on flavoured effects in leptogenesis and lesser people work on nonthermal leptogenesis (cf. [30, 31]). But we attempt to study all the three aspects of leptogenesis in this paper, which makes our work apparently different from others on this account.

The detailed plan of the paper is as follows. In Section 2, methodology and classification of neutrino mass models for zero θ_{13} are presented. Section 3 gives an overview of leptogenesis. The numerical and analytic results for neutrino mass models m_{LL} without and with θ_{13} are given in Sections 4 and 5, respectively. We end with conclusions in Section 6.

2. Methodology and Classification of Neutrino Mass Models

We begin with Type-I seesaw mechanism for estimation of BAU. The required left-handed light neutrino mass models m_{LL} without θ_{13} are given in Table 4. And m_{LL} can be related to the right-handed Majorana mass matrix M_{RR} and the Dirac mass matrix m_{LR} through the inversion seesaw mechanism:

$$M_{RR} = -m_{LR}^{T} m_{LL}^{-1} m_{LR}, \tag{1}$$

where

$$m_{LR} = \operatorname{diag}\left(\lambda^{m}, \lambda^{n}, 1\right) v. \tag{2}$$

In (2) (m, n) are two integers depending on the type of Dirac mass matrix we choose. Since the texture of Yukawa matrix for Dirac neutrino is not known, we take the diagonal texture of m_{LR} to be of charged lepton mass matrix (6, 2), up-quark type mass matrix (8, 4), or down-quark type mass matrix (4, 2), as allowed by SO(10) GUT models.

For computations of leptogenesis, we choose a basis U_R where $M_{RR}^{\text{diag}} = U_R^T M_{RR} U_R = \operatorname{diag}(M_1, M_2, M_3)$ with real and positive eigenvalues. And the Dirac mass matrix m_{LR} in the prime basis transforms to $m_{LR} \rightarrow m'_{LR} = m_{LR} U_R Q$, where Q is the complex matrix containing CP-violating Majorana phases ϕ_1 and ϕ_2 derived from M_{RR}. The values of ϕ_1 and ϕ_2 are chosen arbitrarily other than $\pi/2$ and 0. We then set the Wolfenstein parameter as $\lambda = 0.3$ and compute the three choices of (m, n) in m_{LR}. In this prime basis the Dirac neutrino Yukawa coupling becomes $h = m'_{LR}/v$ and subsequently this value is used in the expression of CP asymmetry. The new Yukawa coupling matrix h also becomes complex, and hence the term $\operatorname{Im}(h^{\dagger}h)_{1j}$ appearing in CP asymmetry parameter ϵ_1 gives a nonzero contribution.

In the second part of this paper, we construct m_{LL} from U_{PMNS} matrix with θ_{13} value:

$$m_{LL} = U_{\text{PMNS}} \cdot m_{\text{diag}} \cdot U_{\text{PMNS}}^{T}, \tag{3}$$

where U_{PMNS} is the Pontecorvo-Maki-Nakagawa-Sakata parameterised matrix taken from the standard particle data group (PDG) [50], and the corresponding mixing angles are

$$\sin^2\theta_{13} = \left|U_{e3}\right|^2, \qquad \tan^2\theta_{12} = \frac{\left|U_{e2}\right|^2}{\left|U_{e1}\right|^2}, \qquad \tan^2\theta_{23} = \frac{\left|U_{\tau 3}\right|^2}{\left|U_{\mu 3}\right|^2}, \tag{4}$$

$$m_{\text{diag}} = \begin{pmatrix} m_1 & 0 & 0 \\ 0 & \pm m_2 & 0 \\ 0 & 0 & m_3 \end{pmatrix}. \tag{5}$$

A global analysis [51, 52] current best-fit data is used in the present analysis:

$$\begin{gathered} \Delta m_{21}^2 = 7.6 \times 10^{-5}\ \text{eV}^2, \qquad \Delta M_{31}^2 = 2.4 \times 10^{-3}\ \text{eV}^2, \\ \sin^2\theta_{12} = 0.312, \qquad \sin^2\theta_{23} = 0.42, \qquad \sin^2\theta_{13} = 0.025, \\ \theta_{12} = 34^\circ \pm 1^\circ, \qquad \theta_{23} = 40.4^{+4.6^\circ}_{-1.8^\circ}, \qquad \theta_{13} = 9.0^\circ \pm 1.3^\circ. \end{gathered} \tag{6}$$

Neutrino oscillation data are insensitive to the low-energy individual neutrino masses. However, it can be measured in tritium beta decay [53] and neutrinoless double beta decay [54] and from the contribution of neutrinos to the energy density of the universe [55]. Very recent data from the Planck experiment have set an upper bound over the sum of all the neutrino mass eigenvalues of $\sum_{i=1}^{3} m_i \leq 0.23\,\text{eV}$ at 95% C.L. [56]. But, oscillations experiments are capable of measuring the two independent mass-squared differences $\Delta m_{21}^2 = m_2^2 - m_1^2$ and $\Delta m_{31}^2 = m_3^2 - m_1^2$ only. This two flavours oscillation approach has been quite successful in measuring the solar and atmospheric neutrino parameters. In the future the neutrino experiments must involve probing the full three flavor effects, including the subleading ones proportional to $\alpha = \Delta m_{21}^2/|\Delta m_{31}^2|$. The Δm_{21}^2 is positive as is required to be positive by the observed energy dependence of the electron neutrino survival probability in solar neutrinos but Δm_{31}^2 is allowed to be either positive or negative by the present data. Hence, two patterns of neutrino masses are possible: $m_1 < m_2 \ll m_3$ called normal hierarchy (NH) where Δm_{31}^2 is positive and $m_3 \ll m_1 < m_2$ called inverted hierarchy (IH) where Δm_{31}^2 is negative. A third possibility, where the three masses are nearly quasi-degenerate with very tiny differences, $m_1 \leq m_2 \sim m_3$, between them, also exists with two subcases of Δm_{31}^2 being positive or negative.

Leptonic CP violation (LCPV) can be established if CP-violating phase δ_{CP} is shown to differ from 0 to 180°. A detailed review on LCPV can be found in [57]. It was not possible to observe a signal for CP violation in the present data so far. Thus, δ_{CP} can have any value in the range $[-180^\circ, 180^\circ]$. The Majorana phases ϕ_1 and ϕ_2 are free parameters. In the absence of constraints on the phases ϕ_1 and ϕ_2, these have been given full variation between 0 and 2π excluding these two extreme values.

3. Leptogenesis

As pointed out above leptogenesis can be thermal or nonthermal; again thermal leptogenesis can be unflavoured (single flavoured) or flavoured which are all explained in the subsequent pages. In the simplest form of leptogenesis the heavy Majorana neutrinos are produced by thermal processes, which is therefore called the "thermal leptogenesis." For our estimations of CP asymmetry parameter ϵ_1 [6, 58, 59], we list here only the required equations for computations. However, interested reader can find more details in [60]. The low-energy neutrino physics is related to the high-energy leptogenesis physics through the seesaw mechanism. In (1), m_{LR}^T is the transpose of m_{LR} and m_{LL}^{-1} is the inverse of m_{LL}. For the third generation Yukawa coupling unification, in SO (10) grand unified theory, one obtains the heavy and light neutrino masses as $M_3 \sim \Lambda_{\text{GUT}} \sim 10^{15}\,\text{GeV}$ and $m_3 \sim v^2/M_3 \sim 0.01\,\text{eV}$ respectively. Remarkably, the light neutrino mass m_3 is compatible with $(\Delta m_{\text{atm}}^2)^{1/2} \equiv m_{\text{atm}} \simeq 0.05\,\text{eV}$, as measured in atmospheric neutrino oscillations. This suggests that neutrino physics probes the mass scale of grand unification (GUT), although other interpretations of neutrino masses are possible as well. The heavy Majorana neutrinos have no gauge interactions. Hence, in the early universe, they can easily be out of thermal equilibrium.

This makes the lightest (N_1) of the heavy right-handed Majorana neutrino an ideal candidate for baryogenesis, satisfying the third condition of Sarkarov, the deviation from thermal equilibrium. Assuming hierarchical heavy neutrino masses ($M_1 \ll M_2, M_3$), the CP asymmetry generated due to CP-violating out-of-equilibrium decay of N_1 is given by

$$\epsilon_1 = \frac{\Gamma\left(N_R \longrightarrow l_L + \phi\right) - \Gamma\left(N_R \longrightarrow \bar{l}_L + \phi^{\dagger}\right)}{\Gamma\left(N_R \longrightarrow l_L + \phi\right) + \Gamma\left(N_R \longrightarrow \bar{l}_L + \phi^{\dagger}\right)}, \quad (7)$$

where $\bar{l}_L$ is the antilepton of lepton l_L and ϕ is the Higgs doublets chiral supermultiplets. Consider

$$\epsilon_1 = \frac{3}{16\pi}\left[\frac{\mathrm{Im}\left[\left(h^{\dagger}h\right)_{12}^2\right]}{\left(h^{\dagger}h\right)_{11}}\frac{M_1}{M_2} + \frac{\mathrm{Im}\left[\left(h^{\dagger}h\right)_{13}^2\right]}{\left(h^{\dagger}h\right)_{11}}\frac{M_1}{M_3}\right], \quad (8)$$

where $h = m_{LR}/v$ is the Yukawa coupling of the Dirac neutrino mass matrix in the diagonal basis of M_{RR} and $v = 174\,\mathrm{GeV}$ is the vev of the standard model. At high temperatures, between the critical temperature T_{EW} of the electroweak phase transition and a maximal temperature T_{SPH},

$$T_{EW} \sim 100\,\mathrm{GeV} < T < T_{\mathrm{SPH}} \sim 10^{12}\,\mathrm{GeV}, \quad (9)$$

these processes are believed to be in thermal equilibrium [9]. Although this important phenomenon is accepted by theorists as a correct explanation of baryogenesis via leptogenesis, it is yet to be tested experimentally. Therefore it is very fascinating that the corresponding phenomenon of chirality changing processes in strong interactions might be observed in heavy decay ion collisions at the LHC [61, 62]. The evolution of lepton number (L) and baryon number (B) is given by a set of coupled equations [63] by the electroweak sphaleron processes which violates ($B+L$) but conserves ($B-L$). At temperature T above the electroweak phase transition temperature T_C, the baryon asymmetry can be expressed in terms of ($B-L$) number density as [64]

$$B\left(T > T_C\right) = \frac{8N_F + 4N_H}{22N_F + 13N_H}\left(B - L\right), \quad (10)$$

where ($B-L$) asymmetry per unit entropy is just the negative of the ratio of lepton density n_L and entropy (s), since the baryon number is conserved in the right-handed Majorana neutrino decays. At T_C, any primodial ($B+L$) will be washed out and (10) can be written as [64, 65]

$$\frac{n_B}{s} \simeq -\frac{8N_F + 4N_H}{22N_F + 13N_H}\frac{n_L}{s}. \quad (11)$$

For standard model (SM) the number of fermion families $N_F = 3$, and the number of Higgs doublets $N_H = 1$; and (11) reduces to

$$\frac{n_B}{s} \simeq -\frac{28}{79}\frac{n_L}{s}. \quad (12)$$

The ratio of baryon to photon is not conserved due to variation of photon density per comoving volume [66] at different epoch of the expanding universe. However, for very slow baryon number B nonconserving interactions, the ratio of baryon to entropy in a comoving volume is conserved. Considering the cosmic ray microwave background temperature $T \simeq 2.3\,\mathrm{K}$, we have $s = 7.04 n_{\gamma}$. Here n_{γ} is a photon number density. And finally the observed baryon asymmetry of the universe [67, 68] for the case of standard model is calculated from

$$\eta_B^{\mathrm{SM}} = \left(\frac{\eta_B}{\eta_{\gamma}}\right)^{\mathrm{SM}} \approx 0.98 \times 10^{-2} \times \kappa_1 \epsilon_1. \quad (13)$$

The efficiency or dilution factor κ_1 describes the washout of the lepton asymmetry due to various lepton number violating processes, which mainly depends on the effective neutrino mass

$$\widetilde{m}_1 = \frac{\left(h^{\dagger}h\right)_{11} v^2}{M_1}, \quad (14)$$

where v is the electroweak vev; $v = 174\,\mathrm{GeV}$. For $10^{-2}\,\mathrm{eV} < \widetilde{m}_1 < 10^{-3}\,\mathrm{eV}$, the washout factor κ_1 can be well approximated by [69]

$$\kappa_1\left(\widetilde{m}_1\right) = 0.3\left[\frac{10^{-3}}{\widetilde{m}_1}\right]\left[\log\frac{\widetilde{m}_1}{10^{-3}}\right]^{-0.6}. \quad (15)$$

We adopt a single expression for κ_1 valid only for the given range of $\widetilde{m}_1$ [69–73]. And the comparison of the effective neutrino mass $\widetilde{m}_1$ with the equilibrium neutrino mass

$$m_* = \frac{8\pi H v^2}{M_1^2} \sim 1.1 \times 10^{-3}\,\mathrm{eV} \quad (16)$$

gives the information whether the system is weak or strong washout regime. For the weak washout regime we have $\widetilde{m}_1 < m_*$ and $M_1 \geq 10^{12}\,\mathrm{GeV}$ whereas for the strong washout regime we have $\widetilde{m}_1 > m_*$ and $M_1 \leq 10^{12}\,\mathrm{GeV}$. However, the strong washout regime appears to be favoured by the present evidence for neutrino masses.

In the flavoured thermal leptogenesis [74–77], we look for enhancement in baryon asymmetry over the single flavour approximation and the equation for CP asymmetry in $N_1 \rightarrow l_{\alpha}\phi$ decay where $\alpha = (e, \mu, \tau)$ becomes

$$\begin{aligned}\varepsilon_{\alpha\alpha} = {} & \frac{1}{8\pi}\frac{1}{\left(h^{\dagger}h\right)_{11}} \\ & \times\left(\sum_{j=2,3}\mathrm{Im}\left[h_{\alpha_1}^{*}\left(h^{\dagger}h\right)_{1j}h_{\alpha j}\right]g\left(x_j\right)\right. \\ & \left. + \sum_{j}\mathrm{Im}\left[h_{\alpha_1}^{*}\left(h^{\dagger}h\right)_{j1}h_{\alpha j}\right]\frac{1}{\left(1 - x_j\right)}\right),\end{aligned} \quad (17)$$

where $x_j = M_j^2/M_i^2$ and $g(x_j) \sim (3/2)(1/\sqrt{x_j})$. The efficiency factor is given by $\kappa = m_*/\widetilde{m}_{\alpha\alpha}$. Here too $m_* = (8\pi H v^2/M_1^2) \sim 1.1\times10^{-3}\,\mathrm{eV}$ and $\widetilde{m}_{\alpha\alpha} = (h_{\alpha1}^{\dagger}h_{\alpha1}/M_1)v^2$. This leads to the BAU:

$$\eta_{3B} = \frac{\eta_B}{\eta_{\gamma}} \sim 10^{-2}\sum_{\alpha}\epsilon_{\alpha\alpha}\kappa_{\alpha} \sim 10^{-2} m_*\sum_{\alpha}\frac{\epsilon_{\alpha\alpha}}{\widetilde{m}_{\alpha\alpha}}. \quad (18)$$

Table 1: Predicted values of the solar and atmospheric neutrino mass-squared differences and mixing angles for $\tan^2\theta_{12} = 0.45$.

Type	Δm^2_{21} (10^{-5} eV^2)	Δm^2_{21} (10^{-3} eV^2)	$\tan^2\theta_{12}$	$\tan^2\theta_{23}$	$\sin\theta_{13}$
(IA)	7.82	2.20	0.45	1.0	0.0
(IB)	7.62	2.49	0.45	1.0	0.0
(IC)	7.62	2.49	0.45	1.0	0.0
(IIA)	7.91	2.35	0.45	1.0	0.0
(IIB)	8.40	2.03	0.45	1.0	0.0
(IIC)	7.53	2.45	0.45	1.0	0.0
(III)	7.61	2.42	0.45	1.0	0.0

For single flavour case, the second term in $\epsilon_{\alpha\alpha}$ vanishes when summed over all flavours. Thus

$$\epsilon_1 \equiv \sum_\alpha \epsilon_{\alpha\alpha} = \frac{1}{8\pi}\frac{1}{(h^\dagger h)_{11}}\sum_j \operatorname{Im}\left[\left(h^\dagger h\right)^2_{lj}\right] g\left(x_j\right); \quad (19)$$

this leads to baryon symmetry:

$$\eta_{1B} \approx 10^{-2} m_* \frac{\epsilon_1}{\widetilde{m}} = 10^{-2}\kappa_1\epsilon_1, \quad (20)$$

where $\epsilon_1 = \sum_\alpha \epsilon_{\alpha\alpha}$ and $\widetilde{m} = \sum_\alpha \widetilde{m}_{\alpha\alpha}$. The conditions of weak or strong washout regime for flavoured leptogenesis are the same as in the case of single favoured/unflavoured leptogenesis, however, with one difference that $\widetilde{m}_1$ is the effective mass due to unflavoured leptogenesis while $\widetilde{m}$ is the resultant effective mass due to contributions of three leptons (flavoured leptogenesis).

In nonthermal leptogenesis [78–83] the right-handed neutrinos N_i $(i = 1, 2, 3)$ with masses (M_1, M_2, M_3) produced through the direct nonthermal decay of the inflaton ϕ interact only with leptons and Higgs through Yukawa couplings. The inflaton decay rate Γ_ϕ is given by [30]

$$\Gamma_\phi = \Gamma\left(\phi \longrightarrow N_i N_i\right) \approx \frac{|\lambda|^2}{4\pi} M_\phi, \quad (21)$$

where M_ϕ is the mass of inflaton ϕ. The reheating temperature (T_R) after inflation is [84]

$$T_R = \left(\frac{44}{2\pi^2 g}\right)^{1/4} \left(\Gamma_\phi M_p\right)^{1/2} \quad (22)$$

and the produced baryon asymmetry of the universe can be calculated by the following relation [85]:

$$Y_B = \frac{n_B}{s} = C Y_L = C\frac{3}{2}\frac{T_R}{M_\phi}\epsilon, \quad (23)$$

where $s = 7.0 n_\gamma$ is related to $Y_B = n_B/S = 8.7 \times 10^{-11}$ in (23). From (23) the connection between T_R and M_ϕ is expressed as

$$T_R = \left(\frac{2Y_B}{3C\epsilon}\right) M_\phi. \quad (24)$$

Two boundary conditions in nonthermal leptogenesis are $M_\phi > 2M_1$ and $T_R \leq 0.01 M_1$. The values of M_1 and ϵ for all neutrino mass models are also used in the calculation of theoretical bounds: $T_R^{\min} < T_R \leq T_R^{\max}$ and $M_\phi^{\min} < M_\phi < M_\phi^{\max}$. Only those models which satisfy these constraints can survive in the nonthermal leptogenesis.

4. Numerical Analysis and Results without θ_{13}

We first begin our numerical analysis for m_{LL} without θ_{13} given in the Appendix. The predicted parameters for $\tan^2\theta_{12} = 0.45$, given in Table 1, are consistent with the global best-fit value. For computations of leptogenesis, we employ the well-known inversion seesaw mechanism as explained in Section 2. Finally the estimated BAU for both unflavoured η_{1B} and flavoured η_{3B} leptogenesis for m_{LL} without θ_{13} is tabulated in Table 2. As expected, we found that there is an enhancement in BAU in the case of flavoured leptogenesis η_{3B} compared to unflavoured η_{1B}. We also observe the sensitivity of BAU predictions on the choice of models without θ_{13} and all but the five models are favourable with good predictions (see Table 2). Streaming lining further, by taking the various constraints into consideration, quasi-degenerate Type-1A, QD-1A (6, 2), and NH-III (8, 4) are competing with each other, which can be tested for discrimination in the next level, the nonthermal leptogenesis.

In case of nonthermal leptogenesis, the lightest right-handed Majorana neutrino mass M_1 and the CP asymmetry parameter ϵ_1 are taken from Table 2 and used in all the neutrino mass models m_{LL} while computing the bounds $T_R^{\min} < T_R \leq T_R^{\max}$ and $M_\phi^{\min} < M_\phi \leq M_\phi^{\max}$ and the computed results are tabulated in Table 3. The baryon asymmetry $Y_B = \eta_B/s$ is taken as input value from WMAP observational data. If we compare these calculations with the predictions of certain inflationary models such as chaotic or natural inflationary model which predicts the inflaton mass to be $M_\phi \sim 10^{13}$ GeV, then from Table 3 the neutrino mass models with (m, n) which are compatible with $M_\phi \sim 10^{13}$ GeV are listed as IA-(4, 2), IIB-(4, 2), III-(4, 2), and III-(6, 2) only. The neutrino mass models with (m, n) should be compatible with $M_\phi \sim (10^{10}–10^{13})$ GeV. Again in order to avoid gravitino problem [84] in supersymmetric models, one has the bound on reheating temperature, $T_R \approx (10^6–10^7)$ Gev. This constraint further streamlines the neutrino mass models and the accepted models are IA-(4, 2), IIB-(4, 2), and III-(6, 2) only.

Furthermore, on examination of the predictions of thermal leptogenesis (Table 2) and nonthermal leptogenesis (Table 3) we found that the estimated results are inconsistent with the two mechanisms of leptogenesis in spite of the fact that they are in agreement with the observation separately. Otherwise for a good model we expect these predictions to

TABLE 2: For zero θ_{13}, the lightest RH Majorana neutrino mass M_1 and values of CP asymmetry and baryon asymmetry for QDN models (IA, IB, and IC), IH models (IIA, IIB), and NH models (III), with $\tan^2\theta_{12} = 0.45$, using neutrino mass matrices given in Table 4. The entry (m, n) in m_{LR} indicates the type of Dirac neutrino mass matrix taken as charged lepton mass matrix (6, 2) or up-quark mass matrix (8, 4), or down-quark mass matrix (4, 2) as explained in the text. IA (6, 2) and III (8, 4) appear to be the best models.

Type	(m, n)	M_1	ϵ_1	η_{1B}	η_{3B}	Status
(IA)	(4, 2)	5.43×10^{10}	1.49×10^{-5}	7.03×10^{-9}	2.16×10^{-8}	√
(IA)	(6, 2)	4.51×10^{8}	1.31×10^{-7}	5.76×10^{-11}	1.34×10^{-10}	√
(IA)	(8, 4)	3.65×10^{6}	1.16×10^{-9}	5.72×10^{-13}	1.19×10^{-12}	✗
(IB)	(4, 2)	5.01×10^{9}	2.56×10^{-14}	7.15×10^{-15}	1.09×10^{-9}	✗
(IB)	(6, 2)	4.05×10^{7}	2.06×10^{-16}	5.76×10^{-20}	8.84×10^{-12}	✗
(IB)	(8, 4)	3.28×10^{5}	1.68×10^{-18}	4.67×10^{-22}	7.16×10^{-14}	✗
(IC)	(4, 2)	5.01×10^{9}	1.85×10^{-13}	5.12×10^{-17}	7.16×10^{-9}	✗
(IC)	(6, 2)	4.05×10^{7}	1.47×10^{-15}	3.77×10^{-29}	5.80×10^{-11}	✗
(IC)	(8, 4)	3.28×10^{5}	1.02×10^{-16}	2.82×10^{-20}	4.34×10^{-12}	✗
(IIA)	(4, 2)	4.02×10^{10}	1.12×10^{-12}	2.49×10^{-15}	7.90×10^{-11}	✗
(IIA)	(6, 2)	3.25×10^{8}	9.00×10^{-15}	2.00×10^{-17}	6.34×10^{-13}	✗
(IIA)	(8, 4)	2.63×10^{6}	7.53×10^{-17}	1.67×10^{-19}	5.35×10^{-15}	✗
(IIB)	(4, 2)	9.76×10^{10}	4.02×10^{-6}	3.25×10^{-9}	7.53×10^{-9}	✗
(IIB)	(6, 2)	8.10×10^{8}	3.33×10^{-8}	2.57×10^{-11}	5.96×10^{-11}	✗
(IIB)	(8, 4)	6.56×10^{6}	2.71×10^{-10}	2.09×10^{-13}	4.86×10^{-13}	✗
(III)	(4, 2)	3.73×10^{12}	3.09×10^{-5}	8.13×10^{-8}	1.85×10^{-6}	✗
(III)	(6, 2)	4.08×10^{11}	3.74×10^{-5}	7.37×10^{-10}	1.62×10^{-9}	√
(III)	(8, 4)	3.31×10^{9}	3.09×10^{-7}	6.06×10^{-11}	1.13×10^{-10}	√

TABLE 3: Theoretical bound on reheating temperature T_R and inflaton masses M_ϕ in nonthermal leptogenesis, for all neutrino mass models with $\tan^2\theta_{12} = 0.45$. Models which are consistent with observations are marked in the status column.

Type	(m, n)	$T_R^{\min} < T_R \le T_R^{\max}$	$M_\phi^{\min} < M_\phi \le M_\phi^{\max}$	Status
(IA)	(4, 2)	$1.2 \times 10^6 < T_R \le 5.4 \times 10^8$	$1.1 \times 10^{11} < M_\phi \le 4.9 \times 10^{13}$	√
(IA)	(6, 2)	$1.1 \times 10^6 < T_R \le 4.5 \times 10^6$	$9.0 \times 10^{8} < M_\phi \le 3.6 \times 10^{10}$	√
(IA)	(8, 4)	$5.1 \times 10^5 < T_R \le 3.6 \times 10^4$	$7.3 \times 10^{6} < M_\phi \le 9.6 \times 10^{6}$	✗
(IB)	(4, 2)	$6.0 \times 10^{13} < T_R \le 5.0 \times 10^7$	$1.0 \times 10^{10} < M_\phi \le 7.4 \times 10^{3}$	✗
(IB)	(6, 2)	$6.4 \times 10^{13} < T_R \le 4.1 \times 10^5$	$8.1 \times 10^{7} < M_\phi \le 0.51 \times 10^{1}$	✗
(IB)	(8, 4)	$6.4 \times 10^{13} < T_R \le 3.3 \times 10^3$	$6.6 \times 10^{5} < M_\phi \le 3.4 \times 10^{-5}$	✗
(IC)	(4, 2)	$8.9 \times 10^{12} < T_R \le 5.0 \times 10^7$	$1.0 \times 10^{10} < M_\phi \le 5.7 \times 10^{4}$	✗
(IC)	(6, 2)	$9.0 \times 10^{12} < T_R \le 4.1 \times 10^6$	$8.1 \times 10^{7} < M_\phi \le 0.36 \times 10^{1}$	✗
(IC)	(8, 4)	$1.1 \times 10^{12} < T_R \le 3.3 \times 10^3$	$6.6 \times 10^{6} < M_\phi \le 1.8 \times 10^{-2}$	✗
(IIA)	(4, 2)	$1.3 \times 10^{13} < T_R \le 5.0 \times 10^8$	$8.0 \times 10^{10} < M_\phi \le 2.8 \times 10^{6}$	✗
(IIA)	(6, 2)	$1.2 \times 10^{13} < T_R \le 4.1 \times 10^6$	$6.5 \times 10^{8} < M_\phi \le 1.8 \times 10^{2}$	✗
(IIA)	(8, 4)	$1.1 \times 10^{14} < T_R \le 3.3 \times 10^4$	$5.3 \times 10^{6} < M_\phi \le 1.8 \times 10^{-2}$	✗
(IIB)	(4, 2)	$8.9 \times 10^{6} < T_R \le 5.0 \times 10^8$	$2.0 \times 10^{12} < M_\phi \le 7.0 \times 10^{13}$	✗
(IIB)	(6, 2)	$8.0 \times 10^{6} < T_R \le 4.1 \times 10^6$	$1.6 \times 10^{11} < M_\phi \le 9.3 \times 10^{9}$	√
(IIB)	(8, 4)	$7.9 \times 10^{6} < T_R \le 3.3 \times 10^4$	$1.3 \times 10^{9} < M_\phi \le 6.3 \times 10^{5}$	✗
(III)	(4, 2)	$4.0 \times 10^{7} < T_R \le 3.7 \times 10^{10}$	$7.5 \times 10^{11} < M_\phi \le 7.0 \times 10^{15}$	✗
(III)	(6, 2)	$3.6 \times 10^{6} < T_R \le 4.1 \times 10^9$	$8.2 \times 10^{11} < M_\phi \le 9.3 \times 10^{14}$	√
(III)	(8, 4)	$3.5 \times 10^{6} < T_R \le 3.3 \times 10^7$	$6.3 \times 10^{9} < M_\phi \le 6.3 \times 10^{10}$	√

be consistent in both frames of leptogenesis. This implies that there is a problem with neutrino mass models without θ_{13}. Next we study neutrino mass models with nonzero θ_{13} and look for consistency in the predictions of two mechanisms of leptogenesis.

5. Numerical Analysis and Results with θ_{13}

In this section, we investigate the effects of inclusion of nonzero θ_{13} (cf. [1, 2]) in the neutrino mass models and predict the cosmological baryon asymmetry. Unlike in Section 4

Table 4

Type	m_{LL}^{diag}	m_{LL}^{0}	$m_{LL} = m_{LL}^{0} + \Delta m_{LL}$
QDIA	$\text{diag}(1, -1, 1)m_0$	$\begin{pmatrix} 0 & \frac{1}{\sqrt{2}} & \frac{1}{\sqrt{2}} \\ \frac{1}{\sqrt{2}} & \frac{1}{2} & \frac{1}{2} \\ \frac{1}{\sqrt{2}} & -\frac{1}{2} & \frac{1}{2} \end{pmatrix} m_0$	$\begin{pmatrix} x-2y & -ax & -ax \\ -ax & \frac{1}{2}-by & \frac{1}{2}-y \\ -ax & \frac{1}{2}-y & \frac{1}{2}-by \end{pmatrix} m_0$
Input $x = 0.66115$, $y = 0.16535$, $m_0 = 0.4$ (for $\tan^2\theta_{12} = 0.45$, $a = 0.868$, $b = 1.025$)			
QDIB	$\text{diag}(1, 1, 1)m_0$	$\begin{pmatrix} 1 & 0 & 0 \\ 0 & 0 & 1 \\ 0 & 1 & 0 \end{pmatrix} m_0$	$\begin{pmatrix} 1-x-2y & ax & ax \\ ax & 1-by & -y \\ ax & -y & 1-by \end{pmatrix} m_0$
Input $x = 8.314 \times 10^{-5}$, $y = 0.00395$, $m_0 = 0.4\,\text{eV}$ ($a = 0.945$, $b = 0.998$)			
QD1C	$\text{diag}(1, 1, -1)m_0$	$\begin{pmatrix} 1 & 0 & 0 \\ 0 & 0 & 1 \\ 0 & 1 & 0 \end{pmatrix} m_0$	$\begin{pmatrix} x-2y & -ax & -ax \\ -ax & -by & 1-y \\ -ax & 1-y & -by \end{pmatrix} m_0$
Input $x = 8.211 \times 10^{-5}$, $y = 0.00395$, $m_0 = 0.4\,\text{eV}$ ($a = 0.945$, $b = 0.998$)			
IH2A	$\text{diag}(1, 1, 0)m_0$	$\begin{pmatrix} 1 & 0 & 0 \\ 0 & \frac{1}{2} & \frac{1}{2} \\ 0 & \frac{1}{2} & \frac{1}{2} \end{pmatrix} m_0$	$\begin{pmatrix} x-2y & -x & -x \\ -x & \frac{1}{2} & \frac{1}{2}-y \\ -x & \frac{1}{2}-y & \frac{1}{2} \end{pmatrix} m_0$
Inverted hierarchy with even CP parity in the first two eigenvalues (IIA), ($m_i = m_1, m_2, m_3$): $(y/x) = 1.0$, $y = 0.005$, $m_0 = 0.045\,\text{eV}$			
IH2B	$\text{diag}(1, 1, 1)m_0$	$\begin{pmatrix} 0 & 1 & 1 \\ 1 & 0 & 0 \\ 1 & 0 & 0 \end{pmatrix} m_0$	$\begin{pmatrix} 1-2y & x & x \\ x & \frac{1}{2} & \frac{1}{2}-y \\ x & \frac{1}{2}-y & \frac{1}{2} \end{pmatrix} m_0$
Inverted hierarchy with odd CP parity in the first two eigenvalues (IIB), ($m_i = m_1, -m_2, m_3$): $(y/x) = 1.0$, $y = 0.6612$, $m_0 = 0.045\,\text{eV}$			
NH3	$\text{diag}(0, 0, 1)m_0$	$\begin{pmatrix} 1 & 0 & 0 \\ 0 & \frac{1}{2} & -\frac{1}{2} \\ 0 & -\frac{1}{2} & \frac{1}{2} \end{pmatrix} m_0$	$\begin{pmatrix} 0 & -x & -x \\ -x & 1-x & 1-y \\ -x & 1+x & 1-x \end{pmatrix} m_0$
Inputs are $(y/x) = 0.0$, $x = 0.146$, $m_0 = 0.028\,\text{eV}$			

analysis, we do not use the particular form of matrices, but we construct the lightest neutrino mass matrix m_{LL} using (3) through (5). On substituting the observational values [86] into U_{PMNS}, we obtain

$$U_{\text{PMNS}} = \begin{pmatrix} 0.81883 & 0.55230 & 0.156434 \\ -0.48711 & 0.52436 & 0.69840 \\ 0.30370 & -0.64807 & 0.69840 \end{pmatrix}. \tag{25}$$

Using (4), this leads to $\sin^2\theta_{13} = 0.0244716$, $\tan^2\theta_{12} = 0.45495$, and $\tan^2\theta_{23} = 1$. Then the m_{diag} of (5) are obtained from the observation data (cf. [51, 52]) ($\Delta m_{12}^2 = m_2^2 - m_1^2 = 7.6 \times 10^{-5}\,\text{eV}^2$, $\Delta m_{23}^2 = m_2^2 - m_3^2 = 2.4 \times 10^{-3}\,\text{eV}^2$) and calculated out for normal and inverted hierarchy patterns. The mass eigenvalues m_i ($i = 1, 2, 3$) can also be taken from [6, 58, 59]. The positive and negative values of m_2 correspond to Type-IA and Type-IB, respectively. Once the matrix m_{LL} is determined the procedure for subsequent calculations is the same as in Section 4.

Here, we give the result of only the best model due to inclusion of reactor mixing angle θ_{13} in predictions of baryon asymmetry, reheating temperature, and inflaton mass (M_ϕ). Undoubtedly, for $\tan^2\theta_{12} = 0.45$, the best model is NH-IA (6, 2) with baryon asymmetry in unflavoured thermal leptogenesis $B_{uf} = 3.313 \times 10^{-12}$, single flavoured approximation $B_{1f} = 8.844 \times 10^{-12}$, and full flavoured $B_{3f} = 2.093 \times 10^{-11}$. If we examine these values, we find that expectedly there is an enhancement in the predictions of baryon asymmetry parameter by a factor of 10 due to inclusion of flavour effects. Similarly in nonthermal leptogenesis, we found that NH-IA is the best model and the predicted results are

$$\begin{aligned} T_R^{\min} &< T_R \leq T_R^{\max}\,(\text{GeV}) \\ &= 7.97 \times 10^3 < T_R \leq 4.486 \times 10^6, \\ M_\phi^{\min} &< M_\phi \leq M_\phi^{\max}\,(\text{GeV}) \\ &= 8.97 \times 10^8 < M_\phi \leq 2.24 \times 10^{11}. \end{aligned} \tag{26}$$

These results show that the neutrino mass models with θ_{13} are consistent in all the three stages of leptogenesis estimations. And normal hierarchy of Type-IA with charged lepton matrix (6, 2) for diagonal form of Drac mass matrix is the most favoured model out of 18 models. And our calculation for all the models either with or without θ_{13} shows that it is strong washout $\widetilde{m}_1$(or $\widetilde{m}$) $> m_*$ and $M_1 \leq 10^{12}$ GeV, the baryon asymmetry is generated at a temperature $T_R(10^6\,\text{GeV}) < M_1(10^9\,\text{GeV})$ for NH-IA model.

6. Conclusions

We have investigated the comparative studies of baryon asymmetry in different neutrino mass models (namely, QDN, IH, and NH) with and without θ_{13} for $\tan^2\theta_{12}$, and we found that models with θ_{13} are better than models without θ_{13}. The predictions of any models with zero θ_{13} are haphazard in spite of the fact that their predictions are consistent in a piecemeal manner with the observational data (see Tables 2 and 3) whereas the predictions of any models with nonzero θ_{13} are consistent throughout the calculations. And among them, only the values of NH-IA (6, 2) satisfied Davidson-Ibarra upper bound on the lightest RH neutrino CP asymmetry $|\epsilon_1| \leq 3.4 \times 10^{-7}$ and M_1 lies within the famous Ibarra-Davidson bound; that is, $M_1 > 4 \times 10^8$ GeV [87]. Neutrino mass models either with or without θ_{13}, Type-IA for charged lepton matrix (6, 2) in normal hierarchy appears to be the best if $Y_B^{\text{CMB}} = 6.1 \times 10^{-10}$ is taken as the standard reference value; on the other hand if then charged lepton matrix (5, 2) is not ruled out. We observed that unlike neutrino mass models with zero θ_{13}, where μ predominates over e and τ contributions, for neutrino mass models with nonzero θ_{13}, τ predominates over e and μ contributions. This implies the factor changes for neutrino mass models with and without θ_{13}. When flavour dynamics is included the lower bound on the reheated temperature is relaxed by a factor ~3 to 10. We also observe enhancement effects in flavoured leptogenesis compared to nonflavoured leptogenesis by one order of magnitude. Such predictions may also help in determining the unknown Dirac Phase δ in lepton sector, which we have not studied in the present paper. And our calculations show that the strong washout regime holds which is favoured by the current evidence for neutrino masses; the baryon asymmetry is generated at a temperature $T_R(10^6\,\text{GeV}) < M_1(10^9\,\text{GeV})$ for NH-IA model. The overall analysis shows that normal hierarchical model appears to be the most favourable choice in nature. Further enhancement from brane world cosmology [88] may marginally modify the present findings, which we have kept for future work.

Appendix

Classification of Neutrino Mass Models with Zero

We list here the zeroth order left-handed Majorana neutrino mass matrices m_{LL}^0 [89–92] with texture zeros left-handed Majorana neutrino mass matrices, $m_{LL} = m_{LL}^0 + \Delta m_{LL}$, corresponding to three models of neutrinos, namely, quasi-degenerate (QD1A, QD1B, and QD1C), inverted hierarchical (IH2A, IH2B), and normal hierarchical (NH3) along with the inputs parameters used in each model. m_{LL} which obey μ-τ symmetry are constructed from their zeroth-order (completely degenerate) mass models m_{LL}^0 by adding a suitable perturbative term Δm_{LL}, having two additional free parameters. All the neutrino mass matrices given in Table 4 predict $\tan^2\theta_{12} = 0.45$. The values of three input parameters are fixed by the predictions on neutrino masses and mixings in Table 1.

Conflict of Interests

The author declares that there is no conflict of interests regarding the publication of this paper.

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