Leptonic CP violation and leptogenesis in minimal supersymmetric $SU(4)_c \times SU(2)_L \times SU(2)_R$

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We consider a supersymmetric $SU(4)_c \times SU(2)_L \times SU(2)_R$ model with a minimal number of Higgs multiplets and Dirac and Majorana CP -violating phases in the neutrino flavor mixing matrix. The model incorporates the charged fermion masses and quark mixings, and uses type I seesaw to explain the solar and atmospheric neutrino oscillations. With the neutrino oscillation data of two mass squared differences and three flavor mixing angles, we employ thermal leptogenesis and the observed baryon asymmetry to find the allowed regions for the Dirac and Majorana phases. For a normal neutrino mass hierarchy, we find that the observed baryon asymmetry can be reproduced by a Dirac phase of around $\delta_{CP} = 3\pi/2$, which is strongly indicated by the recent T2K and NO ν A data. For the case of inverted neutrino mass hierarchy, the predicted baryon asymmetry is not compatible with the observed value.

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The neutrino oscillation phenomena have established non-zero neutrino masses and mixings between different neutrino flavors. Two neutrino mass squared differences and three mixing angles in the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) mixing matrix are measured with good accuracy. The neutrino oscillation parameters to be determined in future experiments include the Dirac CP -violating phase (δ_{CP}) in the PMNS mixing matrix and the ordering of the neutrino mass eigenvalues. The recent results by the T2K experiment [1] strongly indicate a CP violation in the lepton sector with the Dirac CP -violating phase of around $\delta_{CP} = \frac{3\pi}{2}$. The T2K results are also consistent with the results by the NO ν A experiment [2]. In the not too distant future, the accuracy of measurements for the CP -violating phase will be significantly improved. A precise information of quark and lepton mass matrices could provide important clues regarding the origin of fermion masses, flavor mixings and CP violations which, most likely, come from new physics beyond the Standard Model (SM).

In order to explain the observed neutrino masses and flavor mixings, we need to extend the SM. The type I seesaw mechanism [3] is one of the promising ways not only to incorporate the neutrino masses and flavor mixings but to also explain the tiny of neutrino masses naturally. A class of supersymmetric (SUSY) grand unified theories (GUT) has

attracted much interest in this regard. In addition to providing a resolution of the gauge hierarchy problem, the paradigm of SUSY grand unification is also supported by the successful unification of the three SM gauge couplings at the GUT scale, $M_{\text{GUT}} \simeq 2 \times 10^{16}$ GeV. Among several possibilities, SO(10) unification is one of the more compelling ones, with the quark and lepton multiplets of each generation unified in a **16**-dimensional representation along with a right-handed neutrino. The seesaw mechanism is also automatically implemented, being associated with the breaking of SO(10) symmetry to the SM gauge group at M_{GUT} , which is fairly close to the desired seesaw scale.

The so-called minimal SUSY SO(10) model [4] with the minimal set of Higgs multiplets (**10** + **$\overline{126}$**) relevant for fermion mass matrices is a natural extension of nonsupersymmetric SO(10) models considered a long time ago [5]. Because of the unification of quarks and leptons in the **16** representation and the minimal set of Higgs multiplets, the fermion Yukawa matrices are highly constrained with the quark and lepton mass matrices related to each other. Note that the Higgs 10-plet has been used to implement $t - b - \tau$ Yukawa unification in SO(10) [6]. There have been several efforts within the SO(10) framework to simultaneously reproduce the observed quark-lepton mass matrix data as well as the neutrino oscillation data [7–9]. It is quite interesting that after the data fitting, essentially no free parameter is left and all fermion Yukawa matrices, in particular the neutrino Dirac Yukawa matrix, are unambiguously determined. The neutrino Dirac Yukawa matrix allows us to provide concrete predictions for proton lifetime [10] and the rate of lepton flavor violations [11].

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However, the minimal SO(10) model suffers from a serious problem. The observed neutrino oscillation data suggest the right-handed neutrino mass scale to be around 10^{13} – 10^{14} GeV, which is a few orders of magnitude below the GUT scale. With fixed Yukawa couplings of right-handed neutrinos in the minimal SO(10) model, this intermediate scale is provided by the vacuum expectation value (VEV) of the $\overline{126}$ Higgs multiplet. This indicates the existence of many exotic states with intermediate mass scale, which significantly alter the running of the minimal SUSY SM (MSSM) gauge couplings. This has been discussed in Ref. [12], where it is shown that the gauge couplings are not unified any more, and even the SU(2) gauge coupling blows up below the M_{GUT} . To solve this problem, we may extend the minimal model or may consider a different direction in constructing GUT models [13].

In this paper we consider a supersymmetric $SU(4)_c \times SU(2)_L \times SU(2)_R$ (4-2-2) model with a set of Higgs multiplets which closely resembles the minimal SO(10) model. The Higgs multiplets which play an important role in our discussion are $H_{1,2,2}:(\mathbf{1}, \mathbf{2}, \mathbf{2})$, $H_{15,2,2}:(\mathbf{15}, \mathbf{2}, \mathbf{2})$, $H_{10,1,3}:(\mathbf{10}, \mathbf{1}, \mathbf{3})$, and $\bar{H}_{10,1,3}:(\overline{\mathbf{10}}, \mathbf{1}, \mathbf{3})$ corresponding to the Higgs multiplets in the minimal SO(10) model, namely $H_{1,2,2} \subset \mathbf{10}$ -plet Higgs and $H_{15,2,2} + H_{10,1,3} \subset \overline{126}$ -plet Higgs. In the SO(10) model context, the Higgs multiplet $\bar{H}_{10,1,3}$ belongs to 126-plet Higgs, which is introduced to satisfy the D -flat condition along with the $\overline{126}$ -plet Higgs. The $SU(4)_c \times SU(2)_L \times SU(2)_R$ symmetry is broken down to the MSSM gauge group by VEVs of $\langle H_{10,1,3} \rangle = \langle \bar{H}_{10,1,3} \rangle$ satisfying the D -flat condition. Although we are not going into details of the Higgs potential, it is easy to realize a successful Higgs sector of our model by analogy to the minimal set of Higgs multiplets in the minimal renormalizable SO(10) model [14]. In the SO(10) model, the minimal set of Higgs multiplets consists of $\mathbf{10}$ -plet, $\mathbf{126}$ -plet, $\overline{126}$ -plet and $\mathbf{210}$ -plet. It has been demonstrated in Ref. [14] that the most general renormalizable superpotential for the minimal set of Higgs multiplets realizes the SO(10) symmetry breaking to the MSSM gauge group, leaving only the MSSM particle contents light. Since our 4-2-2 model can be embedded in the minimal renormalizable SO(10) model, we can consider a successful Higgs sector of our model as a subset of the Higgs sector of the minimal SO(10) model.

The superpotential relevant for the fermion mass matrices is given by

$$W = Y_1^{ij} F_i \bar{F}_j H_{1,2,2} + Y_{15}^{ij} F_i \bar{F}_j H_{15,2,2} + Y_R^{ij} \bar{F}_i \bar{F}_j H_{10,1,3}, \quad (1)$$

where $F_i:(\mathbf{4}, \mathbf{2}, \mathbf{1})$ and $\bar{F}_i:(\overline{\mathbf{4}}, \mathbf{1}, \mathbf{2})$ denote the matter multiplets in i th generation ($i = 1, 2, 3$). Assuming appropriate VEVs for the Higgs multiplets, we can parametrize the fermion mass matrices as the follows:

$$\begin{aligned} M_u &= c_1 M_{1,2,2} + c_{15} M_{15,2,2}, \\ M_d &= M_{1,2,2} + M_{15,2,2}, \\ M_D &= c_1 M_{1,2,2} - 3c_{15} M_{15,2,2}, \\ M_e &= M_{1,2,2} - 3M_{15,2,2}, \\ M_R &= M_{10,1,3}. \end{aligned} \quad (2)$$

Here M_u, M_d , are the mass matrices for up-type and down-type quarks, M_D is the neutrino Dirac mass matrix, M_e is the charged lepton mass matrix, and M_R is right-handed Majorana neutrino mass matrix. They are given in terms of the three fundamental matrices $M_{1,2,2}, M_{15,2,2}$ and $M_{10,1,3}$ and the complex coefficients c_1 and c_{15} . Note that the relations between fermion mass matrices are exactly the same as those in the minimal SO(10) model, except for M_R . In the 4-2-2 model, M_R is independent of the other mass matrices, while it is proportional to $M_{15,2,2}$ in the minimal SO(10) model.

As is well known, the MSSM gauge couplings successfully unify at $M_{\text{GUT}} \simeq 2 \times 10^{16}$ GeV. In the minimal SO(10) model, M_{GUT} is the scale at which the SO(10) gauge symmetry is broken down to the MSSM gauge symmetry. Since the 4-2-2 model with left-right symmetry¹ can be embedded in the SO(10) model, we simply identify M_{GUT} with the breaking scale of 4-2-2 down to the MSSM gauge group, assuming the left-right symmetry. Therefore, the procedure for fitting the charged fermion mass matrices is the same as in the minimal SO(10) model. On the other hand, it is important to note that M_R being independent of the other mass matrices provides us with the freedom to fit the neutrino oscillation data.

Let us count here the number of free parameters used to fit the charged fermion mass matrices. Because of left-right symmetry, $M_{1,2,2}$ and $M_{15,2,2}$ are 3×3 complex symmetric matrices. Without loss of generality, we take a basis where $M_{1,2,2}$ is real and diagonal, so that the number of free parameters in $M_{1,2,2}$ and $M_{15,2,2}$ is $3 + 12 = 15$. The two complex parameters c_1 and c_{15} introduce an additional 4 degrees of freedom, and therefore in total we have 19 free parameters. The degrees of freedom of charged fermion mass matrices are decomposed into $3 + 6 = 9$ for the lepton and quark mass eigenvalues, and another nine for a unitary matrix for the quark mixings which consists of four parameters in the CKM matrix and five diagonal CP phases. Since the five CP phases are not observable in the SM, we drop these degrees of freedom. Thus, we have 14 free parameters to fit 13 observables [16]. In the minimal SO(10) model, this single free parameter is adjusted to fit the neutrino oscillation data (see [7] for details).

¹The left-right symmetry requires us to add $Y_R^{ij} F_i F_j H_{10,3,1}$ to Eq. (1) with a Higgs multiplet $H_{10,3,1}:(\overline{\mathbf{10}}, \mathbf{3}, \mathbf{1})$. This term corresponds to type II seesaw [15] once $H_{10,3,1}$ develops a nonzero VEV. Since a more complicated Higgs sector seems necessary to induce such a VEV, we do not consider type II seesaw in this paper.

Through the type I seesaw mechanism [3], the light neutrino mass matrix is given by

$$m_\nu = Y_D^T M_R^{-1} Y_D v_u^2 = U_{PMNS}^* D_\nu U_{PMNS}^\dagger, \quad (3)$$

$$U_{PMNS} = \begin{bmatrix} c_{12}c_{13} & c_{12}c_{13} & s_{13}e^{-i\delta_{CP}} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta_{CP}} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta_{CP}} & s_{23}c_{13} \\ s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta_{CP}} & -c_{12}s_{23} - s_{12}s_{23}s_{13}e^{i\delta_{CP}} & c_{23}c_{13} \end{bmatrix} \begin{bmatrix} 1 & 0 & 0 \\ 0 & e^{-i\rho_1} & 0 \\ 0 & 0 & e^{-i\rho_2} \end{bmatrix}, \quad (4)$$

where $c_{ij} = \cos \theta_{ij}$, $s_{ij} = \sin \theta_{ij}$, and ρ_1 and ρ_2 are the Majorana phases. Using Eq. (3), we can express the right-handed neutrino mass matrix as

$$M_R = v_u^2 (Y_D U_{MNS} D_\nu^{-1} U_{MNS}^T Y_D^T). \quad (5)$$

Recall that in this 4-2-2 model, Y_D is fixed by fitting the Dirac fermion masses and mixings in the same manner as the minimal SO(10) model. Hence, employing the current neutrino oscillation data (two neutrino mass squared differences and three mixing angles), we obtain M_R from Eq. (5) as a function of the lightest light neutrino mass eigenvalue, δ_{CP} and $\rho_{1,2}$.

Models with the seesaw mechanism can also account for generating the observed baryon asymmetry in the universe [17],

$$Y_B = \frac{n_B}{s} = (8.6 - 9.0) \times 10^{-11}, \quad (6)$$

via thermal leptogenesis [18], where Y_B is the ratio of the baryon (minus antibaryon) density (n_B) to the entropy density (s). The out-of-equilibrium decays of heavy Majorana neutrinos in the presence of nonzero CP -violating phase generates a lepton asymmetry Y_L in the universe, which is partially converted to the baryon asymmetry

where Y_D is the neutrino Dirac Yukawa matrix and v_u is the VEV of the up-type Higgs doublet. The PMNS mixing matrix, by which m_ν is diagonalized to the mass eigenvalue matrix D_ν , is parametrized as

through (B + L)-violating sphaleron transitions [19,20]. The conversion rate is given by [21]

$$Y_B = -\frac{8N_f + 4N_H}{22N_f + 13N_H} Y_L = -\frac{8}{23} Y_L. \quad (7)$$

Here we set $N_f = 3$ and $N_H = 2$ for the numbers of fermion families N_f and Higgs doublets N_H as in the minimal SUSY SM (MSSM).

The baryon asymmetry produced is evaluated by solving the Boltzmann equations with the information of neutrino Dirac Yukawa coupling matrix and M_R . Since Y_D is fixed and M_R is a function of δ_{CP} and $\rho_{1,2}$ in our model, the resultant baryon asymmetry is given as a function of these parameters. Therefore, leptogenesis constrains the parameters, δ_{CP} and $\rho_{1,2}$, so as to reproduce the observed baryon asymmetry.

As mentioned above, the data fitting procedure for the realistic charged fermion mass matrices is the same as in the minimal SO(10) model, and so in our analysis we employ the numerical values in Y_D found in [7]. In the basis where the charged lepton mass matrix is diagonal, the neutrino Dirac Yukawa coupling matrix at the GUT scale is unambiguously determined and explicitly given by

$$Y_D = \begin{pmatrix} -0.000135 - 0.00273i & 0.00113 + 0.0136i & 0.0339 + 0.0580i \\ 0.00759 + 0.0119i & -0.0270 - 0.00419i & -0.272 - 0.175i \\ -0.0280 + 0.00397i & 0.0635 - 0.0119i & 0.491 - 0.526i \end{pmatrix}, \quad (8)$$

for $\tan \beta = 45$.^{2,3} Using Y_D , we determine M_R from Eq. (5) as a function of δ_{CP} and $\rho_{1,2}$. Since the absolute mass spectrum of light neutrinos has not yet been determined, we consider two cases for it, the normal hierarchical (NH) case and the inverted hierarchical case (IH). For the NH case, the mass eigenvalue matrix D_ν is given by

²Although the output for the neutrino oscillation parameters obtained in [7] is more than 3σ away from the current neutrino oscillation data [22], the experimental data for charged fermion mass matrices are nicely fitted. Since Y_D is determined only by data fitting the charged fermion mass matrices, we can safely use this Y_D data without contradicting any of the experimental results

³In Ref. [23], the charged lepton flavor violating (LFV) processes have been investigated in the minimal SO(10) model with the Y_D data of Eq. (8). Although the right-handed neutrino mass spectrum in our 4-2-2 model is different from the one in the minimal SO(10) model, we expect that the rate of the LFV processes lies in the same order to those presented in [23]. Considering the final results of the MEG experiment [24], $\text{BR}(\mu^+ \rightarrow e^+ \gamma) < 4.2 \times 10^{-13}$ at 90% C.L., we see from the results in [23] that the lower mass bounds on sleptons and winos will be multi-TeV to avoid the MEG constraint.

FIG. 1. For the NH case, heavy neutrino mass spectrum versus m_0 for $\delta_{CP} = \frac{3\pi}{2}$ (left panel), and versus δ_{CP} for $m_0 = 10^{-3}$ eV (right panel). The (green) shaded region denotes the allowed region for δ_{CP} at the 95% confidence level by the recent T2K data [1].

$$D_\nu = \text{diag} \left(m_0, \sqrt{m_0^2 + \Delta m_{12}^2}, \sqrt{m_0^2 + \Delta m_{12}^2 + \Delta m_{23}^2} \right), \quad (9)$$

while for the IH case

$$D_\nu = \text{diag} \left(\sqrt{m_0^2 - \Delta m_{12}^2 + \Delta m_{23}^2}, \sqrt{m_0^2 + \Delta m_{23}^2}, m_0 \right), \quad (10)$$

where the lightest mass eigenvalue m_0 is a free parameter. In our analysis, we adopt the following values for the neutrino oscillation data [22,25]:

$$\begin{aligned} \Delta m_{12}^2 &= 7.6 \times 10^{-5} \text{ eV}^2, \quad \Delta m_{23}^2 = 2.4 \times 10^{-3} \text{ eV}^2, \\ \sin^2(2\theta_{12}) &= 0.87, \quad \sin^2(2\theta_{23}) = 1.0, \quad \sin^2(2\theta_{13}) = 0.092. \end{aligned} \quad (11)$$

Let us first show the mass spectrum of the heavy Majorana neutrinos (mass eigenvalues of M_R) as a function of m_0 and δ_{CP} . For simplicity, we set $\rho_{1,2} = 0$ here. Figure 1 (left panel) shows the mass spectrum M_i ($i = 1, 2, 3$) of the heavy Majorana neutrinos for the NH case as a function of m_0 with $\delta_{CP} = 3\pi/2$ indicated by the recent T2K and

NO ν A data. Since the VEV of $H_{10,1,3}$ (which breaks 4-2-2 down to MSSM) is M_{GUT} , we require $m_0 \gtrsim 10^{-4}$ eV in order to keep Y_R^{ij} within the perturbative regime. The right panel shows the mass spectrum as a function of δ_{CP} for $m_0 = 10^{-3}$ eV. The (green) shaded region denotes the allowed region for δ_{CP} at the 95% confidence level by the recent T2K data [1]. The corresponding results for the IH case are shown in Fig. 2. For the IH case, we find $M_1 \lesssim 10^9$ GeV for any values of m_0 and δ_{CP} .

As shown in Figs. 1 and 2, the heavy neutrino masses are hierarchical for both the NH and IH cases. The lepton asymmetry in the universe in this case is dominantly produced by the lightest heavy neutrino decay, since the asymmetry produced by heavier neutrino decays are almost completely washed out [26]. Thus, we consider the lepton asymmetry produced by only the lightest heavy neutrino decay. In addition, there is a lower bound on the lightest heavy neutrino mass, $M_1 \gtrsim 10^{9-10}$ GeV, to produce the desired amount of baryon asymmetry [27]. For the IH case, the lightest heavy neutrino mass is always found to be below this bound, and in our numerical analysis we find that the

FIG. 2. Same as Fig. 1 but for the IH case.

resultant baryon asymmetry is too small in comparison to the observed baryon asymmetry. Therefore, in the following, we present our results only for the NH case.

For a successful thermal leptogenesis, the reheating temperature (T_r) after inflation must be higher than M_1 for the lightest heavy neutrino to be in thermal equilibrium at T_r . The heavy neutrino mass spectrum shown in Fig. 1 indicates a lower bound on $T_r > 10^9\text{--}10^{10}$ GeV. On the other hand, in SUSY scenarios there is an upper bound on reheating temperature from the cosmological gravitino problem. According to the analysis in Ref. [28], we find $T_r < 10^6\text{--}10^{10}$ TeV depending on the sparticle mass spectrum, in particular with a gravitino mass in the range of $m_{3/2} = 1\text{--}10$ TeV. In our scenario, we assume the gravitino mass of order 10 TeV or higher, so that the reheating temperature can be higher than M_1 while avoiding the cosmological gravitino problem.

Since our model is supersymmetric, we need to consider the lepton asymmetry generated by the decays of both the lightest heavy neutrino and sneutrino. From Fig. 1, the lightest heavy neutrino mass is far below M_{GUT} , and so the effective theory for leptogenesis contains the MSSM and three light neutrinos, as well as the lightest heavy neutrino superfield. Although the complete Boltzmann equations for this system are quite involved (see [29] for complete formulas), because of supersymmetry the lepton asymmetry stored in the SM particles is exactly the same as that stored in the sparticles [29]. Since the heavy neutrino mass scale is much higher than the typical sparticle mass scale $\sim \text{TeV}$, our system is supersymmetric to a very good approximation. Among the many decay and scattering processes involved in the Boltzmann equations, it is known that the (inverse) decay process of the lightest heavy (s) neutrino plays the most important role in determining the resultant baryon asymmetry, while the others are negligible in most of the parameter space [26]. Including only the decay process greatly simplifies the Boltzmann equations, so that for the heavy neutrino they are exactly the same as in the nonsupersymmetric case:

$$\begin{aligned} \frac{dY_{N_1}}{dz} &= \frac{-z}{sH(M_1)} \left(\frac{Y_{N_1}}{Y_{N_1}^{eq}} - 1 \right) \gamma_{N_1}, \\ \frac{dY_{L_f}}{dz} &= -\frac{z}{sH(M_1)} \left[\frac{1}{2} \frac{Y_{L_f}}{Y_{L_f}^{eq}} + \epsilon_1 \left(\frac{Y_{N_1}}{Y_{N_1}^{eq}} - 1 \right) \right] \gamma_{N_1}, \end{aligned} \quad (12)$$

where Y_{N_1} is the yield (the ratio of the number density to the entropy density s) of the lightest heavy neutrino, $Y_{N_1}^{eq}$ is the yield in thermal equilibrium, the temperature of the universe is normalized by the mass of the heavy neutrino $z = M_1/T$, $H(M_1)$ is the Hubble parameter at $T = M_1$, and Y_{L_f} is lepton asymmetry stored in the SM particles. The CP -asymmetry parameter, ϵ_1 , is given by [30]

$$\epsilon_1 = -\frac{1}{2\pi(Y_\nu Y_\nu^\dagger)_{11}} \sum_{j \neq 1} \text{Im}[(Y_\nu Y_\nu^\dagger)_{1j}^2] f(M_j^2/M_1^2), \quad (13)$$

where

$$f(x) \equiv \sqrt{x} \ln \left(\frac{1+x}{x} \right) + 2 \frac{\sqrt{x}}{x-1}, \quad (14)$$

and Y_ν is the neutrino Dirac Yukawa coupling matrix in the basis where both the charged lepton matrix and M_R are diagonalized. Using Eqs. (5) and (8), we can obtain Y_ν as a function of m_0 , δ_{CP} and $\rho_{1,2}$.

The space-time density of the heavy neutrino decay in thermal equilibrium, γ_{N_1} is given by

$$\gamma_{N_1} = s Y_{N_1}^{eq} \frac{K_1(z)}{K_2(z)} \Gamma_{N_1}, \quad (15)$$

where K_1 and K_2 are the modified Bessel functions, and

$$\Gamma_{N_1} = \frac{(Y_\nu Y_\nu^\dagger)_{11}}{8\pi} M_1 \quad (16)$$

is the decay width of the heavy neutrino. Then, we solve the Boltzmann equations with the boundary conditions $Y_{N_1}(0) = Y_{N_1}^{eq}(0)$ and $Y_{L_f}(0) = 0$. The lepton asymmetry generated by the right-handed neutrino decays is converted to the baryon asymmetry via the sphaleron process with the rate of Eq. (7) and hence, we evaluate the resultant baryon number as

$$Y_B = -\frac{8}{23} Y_{L_f}(\infty) \times 2, \quad (17)$$

where the factor 2 takes into account the baryon number stored in sparticles.

FIG. 3. Baryon asymmetry as a function of δ_{CP} for $m_0 = 10^{-3}$ eV, $\rho_2 = 0$ and $\rho_1 = 0$ (solid), $\frac{\pi}{6}$ (dashed) and $\frac{\pi}{3}$ (dotted). The dashed horizontal lines show the range of the observed baryon asymmetry in Eq. (6). The (green) shaded region denotes the allowed region for δ_{CP} at the 95% confidence level by the recent T2K data [1].

FIG. 4. Baryon asymmetry as a function of a Majorana phase ρ_1 for $m_0 = 10^{-3}$ eV, $\delta_{CP} = \frac{3\pi}{2}$ and $\rho_2 = 0$. The dashed horizontal lines show the range of the observed baryon asymmetry in Eq. (6).

For various values of the free parameters (m_0 , δ_{CP} and $\rho_{1,2}$), we numerically solve the Boltzmann equations. In our analysis, we fix $m_0 = 10^{-3}$ eV and $\rho_2 = 0$, for simplicity. For $m_0 \lesssim 10^{-2}$ eV, M_1 is almost independent of m_0 , and we find that the results for the generated baryon asymmetry are almost the same. Figure 3 shows the resultant baryon asymmetries as a function of δ_{CP} for three different values of the Majorana phase, namely, $\rho_1 = 0$ (solid), $\frac{\pi}{6}$ (dashed) and $\frac{\pi}{3}$ (dotted), along with the observed value (horizontal lines). The allowed region for δ_{CP} at the 95% confidence level from the recent T2K data is depicted by the (green) shaded region. We have found the parameters in the shaded region to reproduce the observed baryon asymmetry.

Figure 4 shows the baryon asymmetries as a function of ρ_1 for $\delta_{CP} = \frac{3\pi}{2}$, along with the observed value (horizontal lines). A suitable choice of ρ_1 can reproduce the observed baryon asymmetry.

In summary, we have considered a supersymmetric $SU(4)_c \times SU(2)_L \times SU(2)_R$ model with a minimal number of Higgs multiplets and CP -violating phases (δ_{CP} and $\rho_{1,2}$) in the neutrino flavor mixing matrix. The model has the same structure in the Yukawa couplings for the charged fermions as the supersymmetric minimal $SO(10)$ model, so that the neutrino Dirac Yukawa coupling matrix is unambiguously determined by fitting the experimental data for charged fermion mass matrices. Using the type I seesaw formula with the neutrino Dirac Yukawa coupling matrix, the right-handed Majorana neutrino mass matrix is given as a function of m_0 , δ_{CP} and $\rho_{1,2}$. We have employed leptogenesis and the observed baryon asymmetry to identify the allowed parameter regions. Only the NH case for the light neutrino mass spectrum can reproduce the observed baryon asymmetry with a suitable choice of δ_{CP} and $\rho_{1,2}$. We have found that the Dirac CP -violating phase around $\delta_{CP} = \frac{3\pi}{2}$, which is strongly indicated by the recent T2K and NO ν A data, leads to the baryon asymmetry compatible to the observed value. Once δ_{CP} has been more precisely determined, the allowed regions for $\rho_{1,2}$ will be determined.

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