Split neutrinos, two Majorana and one Dirac, and implications for leptogenesis, dark matter, and inflation

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We propose a simple framework to split neutrinos with a slight departure from tribimaximal—where two of the neutrinos are Majorana type which provide thermal leptogenesis. We propose a model based on S_3 flavor symmetry. The Dirac neutrino with a tiny Yukawa coupling explains primordial inflation and the cosmic microwave background radiation, where the inflaton is the *gauge invariant* flat direction. The observed baryon asymmetry, and the scale of inflation are intimately tied to the observed reactor angle $\sin\theta_{13}$, which can be further constrained by the LHC and the $0\nu\beta\beta$ experiments. The model also provides the lightest right-handed sneutrino as a part of the inflaton to be the dark matter candidate.

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I. INTRODUCTION

It is important to connect the origins of inflation, observed neutrino masses, matter-antimatter/baryon asymmetry, and dark matter within a falsifiable framework of particle physics beyond the standard model (SM), which can be constrained by various low energy observations [1]. Since inflation dilutes all matter except the quantum fluctuations that we see in the cosmic microwave background (CMB) radiation, it is important that the inflaton itself cannot be an arbitrary field; its decay must produce the baryons and the dark matter [2].

Furthermore, in order to explain the observed neutrino masses, one must go beyond the SM. In the simplest setting it is possible to augment the SM gauge group by an extra $U(1)_{B-L}$, whose breaking might be responsible for generating the observed neutrino masses. In a supersymmetric (SUSY) setup, this could be realizable within MSSM $\times$ $U(1)_{B-L}$, where (MSSM stands for the minimal supersymmetric SM).

Gauging $U(1)_{B-L}$ in the SUSY context provides a unique D-flat direction that can be the inflaton candidate as studied previously in Refs. [3,4]. It was pointed out that a small Dirac Yukawa coupling of order $\mathcal{O}(10^{-12})$ can actually help maintain the flatness of the inflationary potential and provide the right amplitude for the density perturbations, and furthermore, the lightest of the right-handed sneutrino (which is now part of the inflaton) can be an excellent dark matter candidate [4].

However, neither inflation nor dark matter requires all the three generations to be Dirac in nature. In fact it is quite plausible that two of the other neutrinos could be Majorana type [5,6]. Since Dirac neutrino mass terms do not carry any lepton number, but the Majorana neutrinos mass terms do, it is now an interesting possibility to realize leptogenesis within this simplest setup.

In this paper we will demonstrate that it is possible to split neutrinos with one Dirac and two Majorana types with a nonvanishing reactor angle $\sin\theta_{13}$, which can explain the baryon asymmetry. The overall neutrino masses are now governed by the Dirac Yukawa and the scale at which the $U(1)_{B-L}$ is broken, therefore achieving inflation, dark matter candidate, neutrino masses, and baryon asymmetry within a common setup. The emphasis will be given on the construction of two Majorana neutrinos with this setup.

II. MOTIVATION FROM INFLATION, DARK MATTER, AND BARYOGENESIS

A. Gauge invariant inflaton

Let us first consider for simplicity a single generation of neutrino with a tiny Dirac Yukawa coupling, h. The superpotential will be given by[1]

$$W \supset hNH^uL, \tag{1}$$

where N, L, and H^u are superfields containing the right-handed neutrino, left-handed lepton, and the Higgs boson that give mass to the up-type quarks, respectively. The above superpotential generates a renoramlizable potential:

$$V(|\sigma|) = \frac{m_\sigma^2}{2}|\sigma|^2 + \frac{h^2}{12}|\sigma|^4 - \frac{Ah}{6\sqrt{3}}|\sigma|^3, \tag{2}$$

[1]The super-Planckian vacuum expectation value (VEV) chaotic inflation based on a right-handed sneutrino as an inflation has been considered in Refs. [7], where the inflaton is an absolute gauge singlet. In our case the inflaton is charged under MSSM $\times$ $U(1)_{B-L}$; therefore, its couplings to the MSSM and $U(1)_{B-L}$ fields are determined by that of the SM and gauge interaction of $U(1)$'s. The advantage of a *gauged inflaton* is that the thermal history of the Universe is under much better control; see [2,8].

where $m_\sigma^2 = (m_{\tilde N}^2 + m_{\tilde L_1}^2 + m_{H^u}^2)/3$ is the soft SUSY mass and it can be in a wide range, i.e. $m_\sigma \geq \mathcal{O}(500)$ GeV, compatible with the current searches of SUSY particles at the LHC. The A term is proportional to the inflaton mass m_σ, and the flat direction field σ is

$$\sigma = (\tilde N + H^u + \tilde L_1)/\sqrt{3}, \tag{3}$$

where $\tilde N$, $\tilde L_1$, H_u are the scalar components of corresponding superfields. Since the right-handed sneutrino $\tilde N$ is a singlet under the SM gauge group, its mass receives the smallest contribution from quantum corrections due to SM gauge interactions, and hence it can be set to be the lightest supersymmetric particle. Therefore the dark matter candidate arises from the right-handed sneutrino component of the inflaton.

Inflation happens near the *inflection point* [9,10],[2] where $A \sim 4m_\sigma$. Near the inflection point it is possible to probe the properties of the inflaton [4]:

$$\sigma_0 \approx \sqrt{3}m_\sigma/h = 6\times 10^{12} m_\sigma(0.05\ \text{eV}/m_\nu), \tag{4}$$

$$V(\sigma_0) \approx (m_\sigma^4/4h^2) = 3\times 10^{24} m_\sigma^4 (0.05\ \text{eV}/m_\nu)^2, \tag{5}$$

where m_ν denotes the neutrino mass that is given by $m_\nu = h\langle H_u\rangle$, with $\langle H_u\rangle \equiv v_u \simeq 174$ GeV. The largest neutrino mass is $m_\nu \sim \mathcal{O}(1)$ eV [12]. The above potential, see Eq. (2), has been studied extensively in order to match the current temperature anisotropy in the CMB radiation. It is possible to match the central values of the temperature anisotropy denoted by $\delta_H = 1.91\times 10^{-5}$ and the spectral tilt: $n_s = 0.968$, see [10], for a wide range of masses $100\ \text{GeV} \leq m_\sigma \leq 10^9$ GeV and the Yukawa for the Dirac neutrino in the range $10^{-12} \leq h \leq 10^{-8}$.

B. Reheat temperature and dark matter abundance

Since the inflaton is a gauge invariant flat direction of MSSM $\times\, U(1)_{B-L}$, it naturally couples only to the MSSM degrees of freedom and the degrees of freedom of the $U(1)_{B_L}$. Note that the $U(1)_{B-L}$ is gauged, so all the inflaton components have gauge interactions. The gauge interaction of $U(1)_{B-L}$ could be set similar to that of the hypercharge, i.e. $g_Y \approx g'_{B-L}$ without any loss of generality. For a certain choice of $g'_{B-L} \sim 0.4$, all the gauge couplings can unify at high scales; see for instance Ref. [4(b)].

The minimum of the inflaton potential, $\sigma = 0$, is a point of *enhanced gauge symmetry* where the entire MSSM $\times$ $U(1)_{B-L}$ gauge symmetry is restored. The gauge bosons, gauginos, and the corresponding Z' and its superpartner coupled to the inflaton become massless at $\sigma = 0$, and they can be excited instantly, nonperturbatively via time-dependent quantum fluctuations, as shown in a detailed analysis of Ref. [8]. The process of thermalization happens within one Hubble time and during this period all the inflaton decay products thermalize efficiently with a reheat temperature governed by the total potential, $V(\sigma_0)$, see Eq. (5), [4]

$$T_{rh} \sim V(\sigma_0)^{1/4} \sim 10^7\ \text{GeV}, \tag{6}$$

for $m_\sigma \sim 500$ GeV, and $m_\nu \sim 0.1$ eV. Scatterings in a thermal bath with the new $U(1)$ gauge interactions also bring the right-handed sneutrino into thermal equilibrium. Note that part of the inflaton, i.e. its $\tilde N$ component, see Eq. (3), has never decayed as it is the lightest supersymmetric particle. However, its gauge interaction with the Z' leads to its efficient annihilation as shown in earlier papers in Ref. [4].

The relic abundance that matches that of the cosmological observations, $\Omega_{\text{CDM}}h^2 \sim 0.12$ [13], is purely set by thermal freeze-out, which was calculated in Refs. [4] for a wide range of the lightest sneutrino mass, i.e. $100\ \text{GeV} \leq m_{\tilde N} \leq 2000$ GeV. The parameter space for which the relic dark matter abundance is matched is found to be quite large for $m_{h^0} > 114.4$ GeV [14] and $1\ \text{TeV} \leq m'_Z \leq 2$ TeV. It would be interesting to reanalyze the parameter space with the new Higgs mass window of $116\ \text{GeV} \leq m_{h^0} \leq 125$ GeV [15] and to revisit the indirect searches for the right-handed sneutrino dark matter [16].

The detailed analysis is beyond the reach of this current paper and we will leave some of these very interesting issues for future investigation.

C. Baryogenesis and the need for two majorana components

Although after inflation the reheat temperature is sufficiently high enough to realize the electroweak baryogenesis within the MSSM, given the current evidence on the Higgs mass searches at the LHC, it is unlikely that the phase transition would be the first order [17]. Therefore, one would have to rely on other ways of generating the observed matter-antimatter asymmetry.

In our case, the Affleck-Dine baryogenesis would not be realizable even if some of the neutrinos are Majorana. Note that the inflaton is comprised of Eq. (3), where all three fields take the same VEVs. Once NLH^u takes a large VEV, other directions such as LH^u, LLe, udd cannot be lifted at higher VEVs simultaneously. All other directions become massive by virtue of the large inflaton's VEV; see for a review [18]. The *only* plausible scenario would be to realize thermal/nonthermal leptogenesis. However, this would require at least two of the neutrinos to be of Majorana type with masses close to the TeV scale. Therefore, upon reheating all three right-handed (s)neutrino components are present in a thermal bath. The lightest of the three sneutrinos is Dirac and a candidate for dark matter while two of the Majorana types would be used to generate the observed lepton asymmetry.

[2]The initial condition for inflation has been discussed in Refs. [11].

III. CONSTRUCTING DIRAC AND MAJORANA NEUTRINOS

Let us first construct the neutrino masses. It was proposed in Ref. [5] that two right-handed neutrinos, namely, N_e and N_τ, could have Majorana mass terms, while the N_μ right-handed neutrino has no Majorana mass at tree-level,[3] and it is coupled to the left-handed neutrino that forms a Dirac mass term. Since the neutrino has a split nature the authors in Ref. [5] call this scenario schizophrenic. However, the Dirac nature is not protected at a higher level, so a schizophrenic case is equivalent to the quasi-Dirac case [19,20]. The overall neutrino mass scale is governed by the Dirac Yukawa h. Therefore, a lower limit for $0\nu\beta\beta$ is obtained in both normal and inverse neutrino mass hierarchy. Since the second neutrino Majorana mass can be made zero, the limit for $0\nu\beta\beta$ is about a factor of 2 larger than the usual Majorana case and this model can be ruled out very soon by the next generation of experiments [21].

The model in Ref. [5] is based on S_3, i.e. the permutation group of three objects; see for instance [22]. Note that S_3 has three irreducible representations, $\mathbf{1}$, $\mathbf{1}'$, and $\mathbf{2}$, where $\mathbf{1}'$ is the antisymmetric singlet. The relevant product rules are $\mathbf{1}' \times \mathbf{1}' = \mathbf{1}$ and $\mathbf{2} \times \mathbf{2} = \mathbf{1} + \mathbf{1}' + \mathbf{2}$. In the basis where the generators of S_3 are real, the product of two doublets, i.e. $a = (a_1, a_2)$ and $b = (b_1, b_2)$, are given by

$$(a_1 b_1 + a_2 b_2)_{\mathbf{1}} + (a_1 b_2 - a_2 b_1)_{\mathbf{1}'} + \begin{pmatrix} a_1 b_2 + a_2 b_1 \\ a_1 b_1 - a_2 b_2 \end{pmatrix}_{\mathbf{2}}. \tag{7}$$

In order to obtain the neutrinos mass matrix, we extend the SM by introducing three right-handed neutrinos: $N_\mu \sim 1$ singlet of S_3 and $N = (N_e, N_\tau) \sim 2$ doublet of S_3. As in Ref. [5], we assume that the combination L_e, L_μ, and L_τ

$$L_2 = \frac{1}{\sqrt{3}}(L_e + L_\mu + L_\tau) \sim \mathbf{1}, \tag{8}$$

transforms as a singlet of S_3, and that

$$L = \begin{pmatrix} L_1 \\ L_3 \end{pmatrix} = \begin{pmatrix} \frac{1}{\sqrt{2}}(L_\mu - L_\tau) \\ \frac{1}{\sqrt{6}}(-2L_e + L_\mu + L_\tau) \end{pmatrix} \sim \mathbf{2}, \tag{9}$$

transform as a doublet of S_3. Equivalently we assume that the right-handed charged leptons combination

$$l_2^c = \frac{1}{\sqrt{3}}(l_e^c + l_\mu^c + l_\tau^c) \sim \mathbf{1}, \tag{10}$$

transforms as a singlet of S_3, and that

[3] Majorana neutrino mass term $N_\mu N_\mu$ is forbidden by means of an Abelian discrete symmetry.

TABLE I. Matter content of the model.

	L_2	L	l_2^c	l^c	N_μ	N	$H_i^{u,d}$	$\varphi^{u,d}$	$\phi_i^{u,d}$	ξ
S_3	**1**	**2**	**1**	**2**	**1**	**2**	**1**, **1**′, **2**	**2**	**1**, **2**	**2**
Z_2	+	−	+	−	+	−	+	−	+	+
Z_2'	+	+	+	−	+	+	+	−	−	+

$$l^c = \begin{pmatrix} l_1^c \\ l_3^c \end{pmatrix} = \begin{pmatrix} \frac{1}{\sqrt{2}}(l_\mu^c - l_\tau^c) \\ \frac{1}{\sqrt{6}}(-2l_e^c + l_\mu^c + l_\tau^c) \end{pmatrix} \sim \mathbf{2}. \tag{11}$$

We assume two Abelian symmetries $Z_2 \times Z_2'$ under which $L_2 \sim (+,+)$, $L \sim (-,+)$, $l_2^c \sim (+,+)$, $l^c \sim (-,-)$, $N_\mu \sim (+,+)$, and $N \sim (-,+)$. In the scalar sector we assume three sets of $SU_L(2)$ Higgs doublets $H^{u,d}$, $\varphi^{u,d}$, and $\phi^{u,d}$. These three sets are distinguished by means of $Z_2 \times Z_2'$ under which they transform as $H^{u,d} \sim (+,+)$, $\varphi^{u,d} \sim (-,-)$, and $\phi^{u,d} \sim (+,-)$, respectively. The matter content of our model is summarized in Table I.

The scalar Higgs doublets $H_i^{u,d} \equiv \{H_1, H_{1'}, H_2\}^{u,d}$ transform as $\mathbf{1}$, $\mathbf{1}'$, and $\mathbf{2}$ with respect to S_3 where $H_2^{u,d} = (H_a^{u,d}, H_b^{u,d})$. Equivalently the Higgs scalar fields $\phi_i^{u,d} \equiv \{\phi_1^{u,d}, \phi_2^{u,d}\}$ with $\phi_2 = (\phi_a^{u,d}, \phi_b^{u,d})$ and $\varphi^{u,d} = (\varphi_a^{u,d}, \varphi_b^{u,d})$ are doublets of S_3.

The superpotential for the scalar fields is given by

$$W_H = \mu_{H_i}^2 H_i^u H_i^d + \mu_\varphi^2 \varphi^u \varphi^d + \mu_{\phi_i}^2 \phi_i^u \phi_i^d + \mu_\xi^2 \xi^u \xi^d. \tag{12}$$

We note that there is no mixing between different sets of Higgs scalars and their masses are free parameters controlled by the corresponding μ^2 parameters. We assume that one Higgs boson has a mass of about 125 GeV while the other scalars are all heavy above 2 TeV. The S_3 symmetry implies that the Higgs boson, for instance, H_a^u and H_b^u, have degenerate masses since they belong to a doublet of S_3. However, in principle, the degeneration can be broken by taking explicitly S_3 breaking terms in the soft sector. The S_3 symmetry can give an interesting prediction for the Higgs mass spectrum that can be studied at the LHC; see for instance [23].

The superpotentials for the lepton sector are given by

$$\begin{aligned} W_l &= y_1^l L_2 l_2^c H_1^d + y_2^l L_2 (l^c \varphi)_2^d + y_3^l (L l^c)_1 \phi_1^d \\ &\quad + y_4^l (L l^c)_2 \phi_2^d, \\ W_\nu &= h L_2 N_\mu H_1^u + h_s (LN)_1 H_1^u + h_a (LN)_{1'} H_{1'}^u \\ &\quad + h_2 (LN)_2 H_2^u. \end{aligned} \tag{13}$$

After electroweak symmetry breaking S_3 is completely broken, namely, $\langle H_a^{0\alpha} \rangle \neq \langle H_b^{0\alpha} \rangle$, $\langle \phi_a^{0\alpha} \rangle \neq \langle \phi_b^{0\alpha} \rangle$, and $\langle \varphi_a^{0\alpha} \rangle \neq \langle \varphi_b^{0\alpha} \rangle$ where $\alpha = u, d$. Under this assumption it is possible to show that $M_l \cdot M_l^\dagger$ can be hierarchical and approximately diagonal, where M_l is the charged lepton mass matrix. So the lepton mixing arises mainly from the neutrino sector. The Dirac neutrino mass matrix is given by

$$m_D = \begin{pmatrix} -\frac{2}{\sqrt{6}} & \frac{1}{\sqrt{3}} & 0 \\ \frac{1}{\sqrt{6}} & \frac{1}{\sqrt{3}} & \frac{1}{\sqrt{2}} \\ \frac{1}{\sqrt{6}} & \frac{1}{\sqrt{3}} & -\frac{1}{\sqrt{2}} \end{pmatrix} \times \begin{pmatrix} h_s v + h_2 u_b & 0 & 0 \\ 0 & hv & 0 \\ 0 & 0 & h_s v - h_2 u_b \end{pmatrix} + \begin{pmatrix} 0 & 0 & -\frac{2}{\sqrt{6}} \\ \frac{1}{\sqrt{2}} & 0 & \frac{1}{\sqrt{6}} \\ -\frac{1}{\sqrt{2}} & 0 & \frac{1}{\sqrt{6}} \end{pmatrix} \times \begin{pmatrix} h_a v' + h_2 u_a & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & -h_a v' + h_2 u_a \end{pmatrix}, \quad (14)$$

where $\langle H_1^{u0}\rangle = v$, $\langle H_{1'}^{u0}\rangle = v'$ $\langle H_a^{u0}\rangle = u_a$, and $\langle H_b^{u0}\rangle = u_b$. Note that in the limit v', $u_a \to 0$ (or h_a, $h_2 \to 0$) the Dirac neutrino mass matrix is diagonalized on the left by tribimaximal mixing U_{TB} [24].[4] For values of h_a, $h_2 \neq 0$ we have deviation from tribimaximal mixing. In particular, we generate a deviation of the reactor angle from zero in agreement with recent T2K [25] and Double Chooz [26] experiments. Apparently the reactor angle is a free parameter in this model (proportional to h_a, h_2); however, we will show below that it is related to the baryon asymmetry.

Let us now consider the right-handed Majorana neutrinos mass terms. We assume a scalar isosinglet (so coupled only to right-handed neutrino mass terms) $\xi = (\xi_a, \xi_b) \sim 2$ doublet of S_3. The superpotential is given by

$$W_M = M(NN)_1 + y_\Delta (NN)_2 \xi. \quad (15)$$

Since the term $N_\mu N_\mu$ is missing, the second neutrino mass state ν_2 does not take a Majorana mass at tree level and gives rise to a quasi-Dirac neutrino mass. Such a term can be forbidden by means of Abelian symmetries. For instance, in Ref. [5], the $N_\mu N_\mu$ term was missing by means of the extra Z_8 symmetry under which $N_\mu \to \omega^6 N_\mu$ and a new scalar isosinglet $X \to \omega X$ where $\omega^8 = 1$. Then the Dirac coupling of N_μ is given by $L_2 N_\mu H X^2/M_{\text{P}}^2$ where M_{P} is the Planck scale.

We assume $\langle \xi_a \rangle = 0$ and $\langle \xi_b \rangle \neq 0$, then the right-handed neutrino mass M_R is diagonal with masses

$$M_{N_e} = M + \Delta, \qquad M_{N_\tau} = M - \Delta, \quad (16)$$

where the two independent free parameters are, respectively, M the $U(1)_{B-L}$ breaking scale, i.e. $M \geq 2$ TeV and $\Delta = y_\Delta \langle \xi_b \rangle$. In the limit $\Delta \ll M$ the two massive right-handed neutrinos have degenerate masses. It would be now desirable to have a mass splitting between N_e and N_τ, since we would like to create the observed matter-antimatter asymmetry in the Universe.

A light neutrino mass matrix arises from a type-I seesaw mechanism [27], $m_\nu = -m_D M_R^{-1} m_D^T$ where m_D is defined in Eq. (14). We assume $u_a = 0$ in Eq. (14) and in order to simplify the notation we observe that the VEVs v, v', and u_b can be reabsorbed with a redefinition of the Yukawa couplings h_s, h_a, and h_2 like $h_\alpha \to v_u h_\alpha / v_\alpha$ where v_u is the standard model Higgs doublet's VEV.

The light neutrino mass matrix is not diagonal in the tribimaximal basis U_{TB}, and it is given by

$$U_{\text{TB}}^T \cdot m_\nu \cdot U_{\text{TB}} = \begin{pmatrix} \frac{h_a^2}{M_{N_\tau}} + \frac{y_1^2}{M_{N_e}} & 0 & h_a\left(\frac{y_1}{M_{N_e}} + \frac{y_2^2}{M_{N_\tau}}\right) \\ 0 & \frac{h}{v_u} & 0 \\ h_a\left(\frac{y_1}{M_{N_e}} + \frac{y_2^2}{M_{N_\tau}}\right) & 0 & \frac{h_a^2}{M_{N_e}} + \frac{y_2^2}{M_{N_\tau}} \end{pmatrix} v_u^2, \quad (17)$$

where $y_1 = h_2 + h_s$ and $y_2 = h_2 - h_s$ where $y_1, y_2 \gg h$. When $h_a = 0$ the above matrix is diagonal. In general the matrix in Eq. (17) is diagonalized by a rotation in the 1–3 plane $R_{13}(\theta)$. The lepton mixing matrix is given by $V_l = U_{\text{TB}} \cdot R_{13}(\theta)$ and the reactor neutrino mixing angle $(V_l)_{13}$ is given by

$$\sin\theta_{13} = \sqrt{\frac{2}{3}} \sin\theta \approx h_a \sqrt{\frac{2}{3}} \frac{(M_{N_e} y_2 + M_{N_\tau} y_1)}{(M_{N_e} y_2^2 - M_{N_\tau} y_1^2)}. \quad (18)$$

The best fit value [28,29] of the reactor neutrino mixing angle is about $\sin\theta_{13} \sim \mathcal{O}(0.1)$. For a small value of the θ angle, the eigenvalues of the matrix in Eq. (17) are approximately given by

$$m_{\nu 1} \approx \left(\frac{h_a^2}{M_{N_\tau}} + \frac{y_1^2}{M_{N_e}}\right) v_u^2, \qquad m_{\nu 2} = h v_u, \qquad m_{\nu 3} \approx \left(\frac{h_a^2}{M_{N_e}} + \frac{y_2^2}{M_{N_\tau}}\right) v_u^2. \quad (19)$$

From this set of equalities we can see immediately that the absolute scale of the neutrinos is fixed by the parameter h that must be less than 10^{-12} in order to have neutrino mass $\mathcal{O}(0.1)$ eV. Note that from inflation $h \geq 10^{-12}$ for electroweak scale soft SUSY breaking masses, therefore predicting a large absolute neutrino mass scale in our case. We can obtain y_1 and y_2 from the two neutrinos square mass difference Δm_{atm}^2 and Δm_{sol}^2, namely,

$$y_1^2 = -h_a^2 \frac{M_{N_e}}{M_{N_\tau}} + \frac{M_{N_e}}{v_u^2} \sqrt{h^2 v_u^2 - \Delta m_{\text{sol}}^2}, \qquad y_2^2 = -h_a^2 \frac{M_{N_\tau}}{M_{N_e}} - \frac{M_{N_\tau}}{v_u^2} \sqrt{h^2 v_u^2 + \Delta m_{\text{atm}}^2 - \Delta m_{\text{sol}}^2}. \quad (20)$$

[4] Tribimaximal mixing U_{TB} is given by the first matrix in Eq. (14).

The parameter h_a is also related to the reactor angle as it is clear from Eq. (18). In particular, we can write $\sin\theta_{13}$ as a function of M, Δ, h, and h_a.

IV. EXPLANATION OF LEPTON ASYMMETRY AND $\sin(\theta_{13})$

Before we discuss the lepton asymmetry, let us recall the thermal history of the Universe. The reheat temperature is sufficiently high enough to bring all the degrees of freedom of MSSM and that of the $U(1)_{B-L}$ in thermal equilibrium. This would mean that all the extra Higgs bosons of $U(1)_{B-L}$ are also in thermal equilibrium along with the two almost degenerate right-handed Majorana (s)neutrinos, whose mass scale is close to $M \sim$ 1–10 TeV, the scale of $U(1)_{B-L}$ breaking. Note that we have taken the soft SUSY breaking mass term to be around 500 GeV, which sets the inflationary scale; see Eq. (2). For this mass range of M the Higgs bosons of the $U(1)_{B-L}$ will not play any role in the electroweak phase transition or in the CP asymmetry for thermal leptogenesis, which we discuss below.

Let us now estimate the required CP asymmetry for thermal leptogenesis. The asymmetry is calculated by the interference diagrams between tree-level and one-loop diagrams, which give rise to $\epsilon = \sum \epsilon_{\beta\beta}$ [30]

$$\epsilon = \frac{\sum_j \mathrm{Im}[(m_D^\dagger m_D)_{1j}]^2}{8\pi(m_D^\dagger m_D)_{11}} g(x_j) \approx \frac{h_a^2}{2\pi}\frac{M}{\Delta}\sin^2\alpha, \tag{21}$$

where $\epsilon_{\beta\beta}$ is the asymmetry of the decay of the right-handed (s)neutrinos into β-family (s)leptons and Higgs bosons, $x_j = M_j^2/M_1^2$, we have used the approximation (in the SUSY case) $g(1+z) \approx 2z^{-1}$, and α is the phase of h_a. The baryon asymmetry is given by

$$Y_B = \eta_B/\eta_\gamma \approx (\epsilon\eta)/g_{\mathrm{SM},} \tag{22}$$

where $g_{\mathrm{SM}} = 118$ and η is the efficiency factor that measures the washout effect, namely, if the right-handed neutrino decays slow enough, its abundance does not decrease according to the Boltzman equilibrium statistics $\eta_N \propto e^{-M_N/T}$ so that the out-of-equilibrium N decays give the lepton asymmetry. The N decay is slow enough if its lifetime is longer than the inverse expansion rate of the Universe $1/H$ where H is the Hubble constant. Then we have

$$R \equiv \frac{\Gamma_N}{H} \sim \frac{m^*}{m_{\nu_1}}, \tag{23}$$

where m_{ν_1} is the lightest neutrino mass and $m^* = 256\sqrt{g_{\mathrm{SM}}}v^2/(3M_P) = 2.3 \cdot 10^3$ eV (see for instance [31]). We can have two extreme situations: $R \ll 1$ so N decays strongly out of equilibrium and then the efficiency factor is $\eta = 1$, or $R \gg 1$ and the lepton asymmetry is suppressed as $\eta \sim 1/R$. In conclusion the efficiency factor can be given approximately as

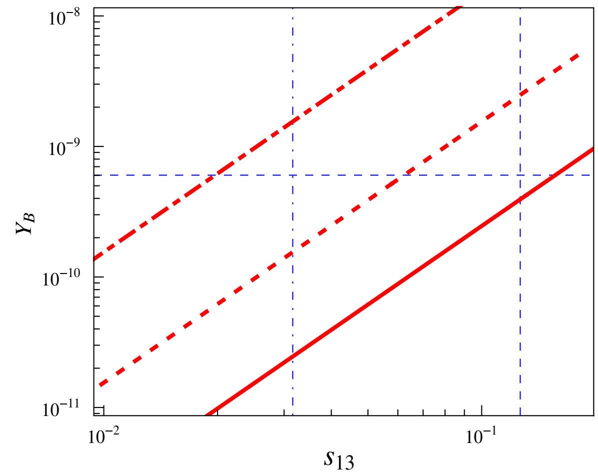

FIG. 1 (color online). Dashed line, $M = 10^3$ GeV, $\Delta = 10^{-6}$ GeV; dotted-dashed line, $M = 10^4$ GeV, $\Delta = 10^{-5}$ GeV; and solid line, $M = 10^3$ GeV, $\Delta = 10^{-5.2}$ GeV fixing $h = 10^{-12}$. The horizontal line is the experimental central value of the baryon asymmetry. The two vertical lines are, respectively, the 3σ lower bound and the best fit values for $\sin\theta_{13}$ [28].

$$\eta \sim \mathrm{Min}(1, m^*/m_{\nu_1}). \tag{24}$$

For fixed values of M, Δ, and α, the baryon asymmetry is a function of the coupling h_a. Equivalently, the reactor angle is a function of h_a if we fix h, $\Delta m^2_{\mathrm{atm}}$, and $\Delta m^2_{\mathrm{sol}}$, besides M, Δ, and α. In Fig. 1, we show the parametric plot of Y_B versus $s_{13} = \sin\theta_{13}$, varying h_a for a different choice of the values of M and Δ by fixing $\Delta m^2_{\mathrm{sol}}$, $\Delta m^2_{\mathrm{atm}}$ at their best fit values, $\alpha = \pi/2$ (maximal CP violation in the lepton sector), and $h = 10^{-12}$ in order to have the neutrino mass scale of about 0.1 eV.

V. SUMMARY

To summarize, we have provided a simple realization of split neutrinos where there is one Dirac neutrino whose light Yukawa coupling explains the flatness of the inflaton potential, the amplitude of the CMB perturbations, and which governs the overall scale of neutrino masses through a $U(1)_{B-L}$ breaking scale. Besides the Dirac neutrino, there are two Majorana neutrinos with a slight departure from tribimaximal mixing, which explains the reactor angle $\sim\theta_{13}$, and tied intimately to the lepton asymmetry obtained from the decay of the two right-handed Majorana (s)neutrinos. This could be a minimal model beyond the SM, where we 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, i.e. the right-handed sneutrino, essentially the inflaton component as a dark matter candidate, and from the $0\nu\beta\beta$ experiments.

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