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Supersymmetric baryogenesis in a hybrid inflation model

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Abstract: We study baryogenesis in a hybrid inflation model which is embedded to the minimal supersymmetric model with right-handed neutrinos. Inflation is induced by a linear combination of the right-handed sneutrinos and its decay reheats the universe. The decay products are stored in conserved numbers, which are transported under the interactions in equilibrium as the temperature drops down. We find that at least a few percent of the initial lepton asymmetry is left under the strong wash-out due to the lighter right-handed (s)neutrinos. To account for the observed baryon number and the active neutrino masses after a successful inflation, the inflaton mass and the Majorana mass scale should be 10^{13} GeV and $\mathcal{O}(10^9$–$10^{10})$ GeV, respectively.

Keywords: Baryo-and Leptogenesis, Cosmology of Theories BSM, Early Universe Particle Physics, Supersymmetry

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1 Introduction

The baryon number of the universe is precisely determined by the observations. For example, the latest result by the Planck collaboration gives [1]

$$(n_B/s)_{\rm obs} = (8.70 \pm 0.04) \times 10^{-11}\,, \tag{1.1}$$

where n_B and s are the number density of the baryon and the entropy density of the present universe, respectively. The genesis of the baryon number is one of the mysteries of our universe since the standard model of particle physics can not explain the observed value of the baryon number. Leptogenesis [2] is a viable mechanism to generate the baryon number. The heavy right-handed neutrinos are introduced in addition to the standard model particles and their decay and scattering in the thermal plasma produce the sufficient lepton density. Finally the lepton number is converted to the baryon number by the sphaleron process. This mechanism is economical in a sense that the heavy right-handed neutrinos also explain the tiny neutrino masses naturally [3–8], called seesaw mechanism.

In this work, we study baryogenesis in a supersymmetric model motivated by the inflation. This model is based on the superconformal subcritical hybrid inflation model [9, 10], where inflation continues even after the inflaton field becomes below the critical point value of the hybrid inflation. Such subcritical regime of inflation is originally considered in refs. [11, 12] with an approximate shift symmetry in Kähler potential.[1] Refs. [9, 10] consider the superconformal model combined with the approximate shift symmetry and found the model gives a good fit with the observed spectral index of the scalar amplitude and the

[1]In the scenario, the waterfall field value is suppressed and inflation continues in the infaton direction, meanwhile refs. [13–15] study the case where inflation along the direction of the waterfall field.

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tensor-to-scalar ratio, and that the inflaton mass is predicted to be around 10^{13} GeV. In the current study, we embed the inflation model to the minimal supersymmetric standard model with right-handed neutrinos. The model is the same as one considered in ref. [16] and we extend the study to a more realistic scenario by taking into account the flavor effects [17–22] and the spectator effects [23–25]. After inflation, the inflaton decays to reheat the universe and produce the lepton numbers. The important points are *i)* not all the lepton numbers are washed out due to the lighter right-handed (s)neutrinos [26, 27] and *ii)* the $B-L$, where B and L are baryon and lepton number respectively, remains due to the conserved charges even in a case where the wash-out effect is most effective [28, 29]. In the latter point the supersymmetry plays the crucial role.

2 The model

We consider an extended minimal supersymmetric standard model (MSSM) augmented by three right-handed neutrinos N_i and two standard model singlet fields $S_\pm$, which are described by a superpotential

$$W \supset \frac{1}{2} M_{ij} N_i N_j + y_{ij} N_i L_j H_u + \lambda_i N_i S_+ S_- \,, \tag{2.1}$$

where y_{ij} and λ_i are coupling constants and M_{ij} are the Majorana masses. The indices i, j take $1, 2, 3$. L_i and H_u are the left-handed lepton doublets and the up-type Higgs, respectively. In this model we assume that $S_\pm$ has a local U(1) charge, $\pm q$, and the other fields are the U(1) singlets, and that the gauge symmetry is spontaneously broken due to the D-term potential, which is studied in ref. [16]. Consequently it acquires nonzero vacuum expectation value (VEV), denoted as $\langle S_+\rangle$. Then in $(N_i, S_-)^T$ basis, we have the following 4 by 4 mass matrix,

$$M_N \equiv \begin{pmatrix} & & & \lambda_1\langle S_+\rangle \\ & M & & \lambda_2\langle S_+\rangle \\ & & & \lambda_3\langle S_+\rangle \\ \lambda_1\langle S_+\rangle & \lambda_2\langle S_+\rangle & \lambda_3\langle S_+\rangle & 0 \end{pmatrix}, \tag{2.2}$$

which leads to a 3 by 3 active neutrino mass matrix,

$$M_\nu = -\langle H_u\rangle^2 \tilde{y}^T M_N^{-1} \tilde{y}, \tag{2.3}$$

$$\tilde{y} = \begin{pmatrix} & y & \\ & & \\ 0 & 0 & 0 \end{pmatrix}. \tag{2.4}$$

This matrix is diagonalized by a unitary matrix U_ν as

$$U_\nu^T M_\nu U_\nu = \mathrm{diag}(m_1, m_2, m_3)\,. \tag{2.5}$$

It is known that the matrix M_ν gives rise to one massless neutrino [16]. We follow the convention such that $m_1 = 0$ and $m_3 > m_2$ for the normal hierarchy (NH) and $m_3 = 0$ and

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$m_2 > m_1$ for the inverted hierarchy (IH). Another important fact is that M_ν is independent of both the mass scale $\lambda_i \langle S_+ \rangle$ and y_{3i}. Therefore, $\lambda_i \langle S_+ \rangle$ and y_{3i} are not constrained by the observed neutrino masses. This is crucial in the later discussion.

For later convenience, we introduce two bases; inflaton basis and mass eigenstate basis. The former one, written as $(N_1', N_2', N_3')^T$, is a basis where N_3' only couples to $S_+ S_-$. Namely $\lambda_i N_i \equiv \tilde{\lambda} N_3'$ and the scalar component of N_3' plays the role of the inflaton field. In $(N_i', S_-)^T$ basis, we have the following 4 by 4 mass matrix,

$$M_N' = \begin{pmatrix} & & & 0 \\ & U_{\rm inf}^* M U_{\rm inf}^\dagger & & 0 \\ & & & m_\phi \\ 0 & 0 & m_\phi & 0 \end{pmatrix}, \tag{2.6}$$

where $N_i' = U_{\inf ij} N_j$ and $m_\phi = \tilde{\lambda} \langle S_+ \rangle$, which corresponds to the inflaton mass. In this basis, $\tilde{y}$ transforms as

$$\tilde{y}' = \begin{pmatrix} & & \\ & U_{\rm inf}^* y & \\ & & \\ 0 & 0 & 0 \end{pmatrix}. \tag{2.7}$$

The latter one is the basis where both M_N' and charged lepton mass matrix are diagonalized. In the basis, the relevant terms in the superpotential are

$$W \supset \frac{1}{2} M_I \hat{N}_I \hat{N}_I + \hat{y}_{Ij} \hat{N}_I \hat{L}_j H_u\,, \tag{2.8}$$

where $\hat{N}_I$ $(I = 1, 2, 3, 4)$ and $\hat{L}_i$ are the mass eigenstates of the heavy right-handed neutrinos plus S_- and the charged leptons.[2] Namely,

$$\hat{N} = U_N^\dagger N'\,, \quad \hat{L} = T^\dagger L\,, \tag{2.9}$$

$$\hat{y} = U_N^T \begin{pmatrix} & & \\ & U_{\rm inf}^* y & \\ & & \\ 0 & 0 & 0 \end{pmatrix} T\,. \tag{2.10}$$

Here we have defined a unitary matrix U_N as

$$U_N^T M_N' U_N = \mathrm{diag}(M_1, M_2, M_3, M_4)\,, \tag{2.11}$$

where $M_1 < M_2 < M_3 < M_4$ and similar for T. Then the PMNS matrix is given by $U_{\rm PMNS} = T^\dagger U_\nu$.

In order not to disturb the inflationary trajectory, we assume

$$\lambda_i \langle S_+ \rangle \gg \Lambda\,, \quad M_{ij} \sim \mathcal{O}(\Lambda)\,, \tag{2.12}$$

[2] We will sometimes use a notation $\hat{L}_\alpha$ where $\alpha = e, \mu, \tau$ in the later discussion.

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where Λ represents the typical scale of the Majorana masses. In this limit, the mass eigenvalues M_I have a relation

$$M_1 \sim M_2 \sim \Lambda\,, \quad m_\phi \simeq M_3 \simeq M_4\,, \tag{2.13}$$

and U_N has a structure as

$$U_N = \left(\begin{array}{c:c} u_{2\times 2} & \mathcal{O}\left(\frac{\Lambda}{m_\phi}\right) \\ \hdashline \mathcal{O}\left(\frac{\Lambda}{m_\phi}\right) & \frac{1}{\sqrt{2}}\begin{pmatrix} i & 1 \\ -i & 1 \end{pmatrix} + \mathcal{O}\left(\frac{\Lambda}{m_\phi}\right) \end{array}\right)\,, \tag{2.14}$$

where $u_{2\times 2}$ is $\mathcal{O}(1)$ 2 by 2 matrix. Eq. (2.12) is the feature of this model and eqs. (2.12)–(2.14) are important in the estimation of the lepton asymmetry. As we described below eq. (2.5), the active neutrino mass matrix is independent of m_ϕ. Thus m_ϕ is not constrained by the observation of the neutrino masses. We will see in the next section that $\hat{y}_{3i}$ are free from the constraint, which means that $\hat{y}_{3i}$ are free parameters in this model.

3 Baryogenesis

The overview of the thermal history of our model is the following:

a) scalar component of N_3', denoted as $\tilde{N}_3'$, drives inflation

b) the inflaton decays to reheat the universe and produce lepton number non-thermally

c) part of the lepton number is washed out by $\hat{N}_1$ and $\hat{N}_2$[3]

d) the lepton number is converted to baryon number by the sphaleron process

At the stage *a)* we can take $\sqrt{2}\,\mathrm{Re}\,\tilde{N}_3' = \phi$ as the inflaton without the loss of generality and ϕ drives inflation. Such inflation models are discussed in refs. [16, 31]. After inflation, the inflaton field decays to reheat the universe. The lepton number is produced simultaneously by the inflaton decay, which is given by

$$L^{\rm dec} \equiv \left.\frac{n_L}{s}\right|_{\rm dec} = \frac{3}{4}\frac{T_R}{m_\phi}\epsilon_\phi\,, \tag{3.1}$$

where T_R and ϵ_ϕ are the reheating temperature and the lepton asymmetry of the inflaton decay, respectively.[4] This corresponds to the stage *b)*. Assuming the instantaneous reheating and $T_R/m_\phi \lesssim 1$, the reheating temperature is given by the decay width Γ_ϕ of the inflaton as $T_R \simeq (90/\pi^2 g_*(T_R))^{1/4}\sqrt{\Gamma_\phi M_{\rm Pl}}$ where $g_*(T_R) \simeq 228.75$ and

$$\Gamma_\phi \simeq \frac{(\hat{y}\hat{y}^\dagger)_{33}}{4\pi} m_\phi\,, \tag{3.2}$$

[3]The decays of $\hat{N}_1$ and $\hat{N}_2$ might give comparable contributions by a tuning of the model parameters [30]. To be conservative we ignore them.

[4]This is similar to right-handed sneutrino inflation and leptogenesis [32–40]. As we will see, however, the thermal history in our model is different from those discussed in the literature.

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where we have used $\Lambda/m_\phi \ll 1$, given in eq. (2.12). Similarly, the asymmetry ϵ_ϕ is given by

$$\begin{aligned}\epsilon_\phi &\simeq \frac{1}{8\pi}\sum_{K\neq 3}\frac{\mathrm{Im}[\{(\hat{y}\hat{y}^\dagger)_{3K}\}^2]}{(\hat{y}\hat{y}^\dagger)_{33}}g(x_K)\\ &\simeq -\frac{1}{8\pi}\frac{\mathrm{Im}[\{(\hat{y}\hat{y}^\dagger)_{34}\}^2]}{(\hat{y}\hat{y}^\dagger)_{33}}\frac{m_\phi}{\Delta M}\,,\end{aligned} \tag{3.3}$$

where $g(x) = (\frac{2}{1-x} - \ln\frac{1+x}{x})\sqrt{x}$, $x_K = M_K^2/M_3^2$ and $\Delta M = M_4 - M_3 \sim \mathcal{O}(\Lambda)$. The second line comes from $x_{1,2} \ll 1$ and $x_4 \simeq 1$. Then $g(x_4) \simeq -M_3/\Delta M \simeq -m_\phi/\Delta M$ gives the dominant contribution. Here we have used the fact

$$\frac{(\Gamma_\phi M_3)^2}{(M_3^2 - M_4^2)^2} \simeq \left[\frac{\Gamma_\phi}{2\Delta M}\right]^2 \sim \left[4\times 10^{-5}\ \frac{(\hat{y}\hat{y}^\dagger)_{33}}{10^{-6}}\frac{m_\phi}{10^{13}\,\mathrm{GeV}}\frac{10^{10}\,\mathrm{GeV}}{\Lambda}\right]^2 \ll 1\,, \tag{3.4}$$

in the parameter space we are interested in.

To estimate T_R and ϵ_ϕ, it is convenient to introduce a 4 by 3 matrix R based on ref. [41]:

$$R = iD_N^{-1/2}U_N^T\tilde{y}'\langle H_u\rangle U_\nu D_\nu^{-1/2}\,, \tag{3.5}$$

where $D_N^{\pm1/2} = \mathrm{diag}(M_1^{\pm1/2}, M_2^{\pm1/2}, M_3^{\pm1/2}, M_4^{\pm1/2})$, and $D_\nu^{\pm1/2} = \mathrm{diag}(0, m_2^{\pm1/2}, m_3^{\pm1/2})$ for the NH, $\mathrm{diag}(m_1^{\pm1/2}, m_2^{\pm1/2}, 0)$ for the IH. R satisfies $R^TR = \mathrm{diag}(0,1,1)$ and $\mathrm{diag}(1,1,0)$ for the NH and IH, respectively. From the equation, we write $\hat{y}$ in terms of R as

$$\hat{y}\langle H_u\rangle = -iD_N^{1/2}RD_\nu^{1/2}U_{\mathrm{PMNS}}^\dagger\,, \tag{3.6}$$

which leads to

$$(\hat{y}\hat{y}^\dagger)\langle H_u\rangle^2 = D_N^{1/2}RD_\nu R^\dagger D_N^{1/2}\,. \tag{3.7}$$

In the expansion of Λ/m_ϕ, we find

$$(\hat{y}\hat{y}^\dagger)_{33} = m_\phi\sum_i m_i|R_{3i}|^2/\langle H_u\rangle^2\,, \tag{3.8}$$

$$(\hat{y}\hat{y}^\dagger)_{34} = i(\hat{y}\hat{y}^\dagger)_{33}(1 + \mathcal{O}(\Lambda/m_\phi))\,, \tag{3.9}$$

where we have used eq. (2.14). It is clear that $\hat{y}_{3i}$ is not constrained by the observed results in the neutrino sector. We note that imaginary part of $\{(\hat{y}\hat{y}^\dagger)_{34}\}^2$ is suppressed by Λ/m_ϕ compared with the naive expectation, which cancel a factor of $m_\phi/\Delta M$. Thus, we get

$$T_R \simeq 6\times 10^{11}\,\mathrm{GeV}\sqrt{\frac{(\hat{y}\hat{y}^\dagger)_{33}}{10^{-6}}\frac{m_\phi}{10^{13}\,\mathrm{GeV}}}\,, \tag{3.10}$$

$$\epsilon_\phi \simeq \frac{a}{4\pi}(\hat{y}\hat{y}^\dagger)_{33} \simeq 8\times 10^{-8}a\frac{(\hat{y}\hat{y}^\dagger)_{33}}{10^{-6}}\,. \tag{3.11}$$

Here we have introduced a coefficient a to take into account $\mathcal{O}(\Lambda/m_\phi)$ term in eq. (3.9). We expect $a = \mathcal{O}(1)$ without a fine-tuning. T_R/m_ϕ can be written in terms of K_ϕ as $T_R/m_\phi \simeq \sqrt{K_\phi}$. Here $K_\phi \equiv K_3 \simeq K_4$, where $K_I \equiv \tilde{m}_I/m_*$ and

$$\tilde{m}_I \equiv \frac{(\hat{y}\hat{y}^\dagger)_{II}\langle H_u\rangle^2}{M_I}\,, \quad m_* = \frac{4\pi^2\sqrt{g_*(M_I)}\langle H_u\rangle^2}{3\sqrt{10}M_{\mathrm{Pl}}}\,. \tag{3.12}$$

Thus the condition $T_R/m_\phi \lesssim 1$ is equivalent to $K_\phi \lesssim 1$.

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At the stage *c)*, the generated lepton number suffers from the wash-out by $\hat{N}_1$ and $\hat{N}_2$. To evaluate the wash-out effect, we estimate $K_{1,2}$. From eq. (3.7), it is straightforward to obtain

$$\tilde{m}_I = \sum_i m_i |R_{Ii}|^2 \geq m_{\rm min}\,, \tag{3.13}$$

for $I = 1, 2$ where $m_{\rm min} = m_2 \simeq 8.6 \times 10^{-3}\,\mathrm{eV}$ and $m_1 \simeq 4.9 \times 10^{-2}\,\mathrm{eV}$ for the NH and IH cases, respectively [42]. Here we have taken $\Lambda/m_\phi \ll 1$ to derive $m_{\rm min}$. Therefore, $K_{1,2}$ have the minimum values as

$$K_I \geq \begin{cases} 22 & (\mathrm{NH}) \\ 124 & (\mathrm{IH}) \end{cases} \qquad \text{for } I = 1, 2\,. \tag{3.14}$$

This means that the wash-out effect is strong. Even in the strong wash-out regime, not all lepton number is washed out [26, 27]. At the production of the lepton number we assume

$$m_\phi \gtrsim 10^{13}\,\mathrm{GeV}\,. \tag{3.15}$$

This requirement is for a successful inflation, which will be quantified later in eq. (3.23). The mass scale of inflaton means that the all Yukawa interactions are out of equilibrium at the decay of the inflaton, except for the top Yukawa interaction. Consequently the produced lepton is a coherent state $|\ell_3\rangle\,(\simeq |\ell_4\rangle) \equiv |\ell_\phi\rangle$, defined by $|\ell_I\rangle = \frac{1}{\sqrt{(\hat{y}\hat{y}^\dagger)_{II}}} \hat{y}_{I\alpha}\,|\ell_\alpha\rangle$. As the temperature drops down, the spectator effects [23–25], the flavor effects [17–22] and the wash-out effect due to $\hat{N}_1$ and $\hat{N}_2$ become important.

Regarding the masses of the lighter right-handed neutrinos, we consider

$$10^7\,\mathrm{GeV} \lesssim M_{1,2} \lesssim 10^{10}\,\mathrm{GeV}\,. \tag{3.16}$$

Here the upper bound is from the requirement (2.12), meanwhile the lower one is to ignore the μ term, i.e., $\mu H_u H_d$, and the gaugino masses. To be more quantitative, $M_{1,2}$ should be larger than roughly $2 \times 10^7\,\mathrm{GeV}(\mu/100\ \mathrm{GeV})^{2/3}$ and $8 \times 10^7\,\mathrm{GeV}(m_{\tilde{g}}/1\ \mathrm{TeV})^{2/3}$ [29, 43]. As temperature gets down to $T \sim M_{1,2}$, the lepton number is transported by the interactions that are in equilibrium. In the MSSM, there are 18 independent fields and 13 types of interactions [44]. Five U(1) charges are anomaly-free, which are hypercharge, $\Delta_\alpha \equiv B/3 - L_\alpha$ ($\alpha = e, \mu, \tau$), and $\mathcal{R}$ defined in [44]. Here B and L_α are the baryon number and the lepton number of each flavor. It is worth noting that $\mathcal{R}$ is different from the R-symmetry of the supersymmetric model. The rest of 13 U(1) charges are broken when the all interactions are in equilibrium. We can take a convenient linear combination of the charges, which are broken one by one when an interaction enters in equilibrium as the temperature gets lower.[5]

[5]We take into account the neutrino Yukawa interactions later. Since the right-handed neutrinos are gauge singlets, their scalar and fermionic components are independent. Therefore, two additional degrees of freedom with two interaction terms, i.e., Majorana masses and the neutrino Yukawas, lead to no additional U(1) and the five anomaly free U(1)s are kept unbroken. Our model has two new fields $S_\pm$ with one Yukawa interaction. Therefore a U(1) should appears, which corresponds to the gauged U(1) in our setup. After inflation, S_+ gets the VEV and the gauged U(1) is broken. At the same time S_+ obtains a mass of $\sqrt{\xi} \sim \mathcal{O}(10^{15})$ GeV, which is integrated out in the energy scale we are interested in. The mass scale of S_- is the order of the inflaton mass. Therefore, it is also integrated out below the energy scale of the reheating, as well as the inflaton field.

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The U(1) charges are listed in appendix A. The crucial point is that at the decay of the inflaton non-zero conserved charges are created in addition to Δ_α. As a consequence, the lepton number stored in the conserved charges escapes from the wash-out by $\hat{N}_1$ and $\hat{N}_2$.[6] On the other hand, Δ_α are affected by the wash-out effect. Therefore, we have two contributions to the $B-L$, the conserved charges and the relic that survives the wash-out.

In the present case, left-handed (s)leptons and up-type Higgsinos (Higgses) are produced by the inflaton decay. Additionally, $\mathcal{R}$ and R^χ_{3q} are expected to be generated. This is due to the R symmetry breaking during the coherent inflaton oscillation. Therefore the inflaton decay gives an initial condition of the chemical potentials of the conserved charges are written as

$$\mu_{Q^i} = (\mu_R, -r_e\mu_{L^{\rm dec}}, -r_\mu\mu_{L^{\rm dec}}, -r_\tau\mu_{L^{\rm dec}}, \mu_R, \mu_R, \mu_R, \mu_R, \mu_R)\,, \tag{3.17}$$

where $Q^i = (\mathcal{R}, \Delta_e, \Delta_\mu, \Delta_\tau, R^\chi_{3u}, R^\chi_{3d}, R^\chi_{3s}, R^\chi_{3c}, R^\chi_{3b})$ and the other chemical potentials are zero. $\mu_{L^{\rm dec}}$ is the chemical potential of the total lepton number produced at the inflaton decay, i.e., $n_L|_{\rm dec} = \mu_{L^{\rm dec}}T^2/6$ and r_α are the fractions of each flavor. μ_R is the chemical potential of R number.

A finite value of μ_R may come from inflation or the coherent oscillation of the inflaton field. In the VEV $\langle\phi\rangle$ or the variance $\sqrt{\langle\phi^2\rangle}$ of the inflaton field induces additional one-loop diagram with R-breaking intermediate states appears, which leads to an asymmetry of R. With the variance, for instance, the asymmetry is estimated to be suppressed by at least $(\hat{y}\hat{y}^\dagger)_{33}\langle\phi^2\rangle/m_\phi^2$, compared with ϵ_ϕ. Here we found the same suppression factor Λ/m_ϕ from the imaginary part of the Yukawa couplings as ϵ_ϕ. Then, it is $\mathcal{O}(10^{-10})$ suppression in our target parameter space.[7] Therefore it can be ignored in our current study and we omit the contribution from μ_R in the discussion below. On the other hand, it may be an importnat contribution to the baryon assymmetry if $(\hat{y}\hat{y}^\dagger)_{33} \sim 1$. In that case, we need to take into account the non-perturbative decay of inflaton during the coherent oscillation. Or if we consider a different type of the seesaw mechanism, the suppression factor Λ/m_ϕ may be irrelevant to boost the asymetric parameter. We will leave these possibilities for the future research.

Let us consider $M_{1,2} \sim 10^9$–10^{10} GeV. Using eq. (B.6) in appendix B, we get

$$\mu_{Q^{B-L}} = \frac{120251}{148420}\mu_{Q^{\Delta_0}} \simeq 0.81\mu_{Q^{\Delta_0}}\,, \tag{3.18}$$

Referring to ref. [27], $\mu_{Q^{\Delta_0}}$ can be obtained as follows. At the temperature the QCD and electroweak sphaleron processes, b, τ, and c Yukawa interactions are in equilibrium. Due to the τ Yukawa interaction, $|\ell_\tau\rangle$ component in $|\ell_\phi\rangle$ is washed out since $K_I^\tau \gg 1$ for $I = 1, 2$, where $K_I^\alpha \equiv K_I|\langle\ell_\alpha|\ell_I\rangle|^2$ $(\alpha = e, \mu, \tau)$.[8] Let us call $|\ell_\tau^\perp\rangle$ as a state orthogonal to $|\ell_\tau\rangle$. Next, we decompose $|\ell_\tau^\perp\rangle$ by states $|\ell_I^{\tau\perp}\rangle$ and $|\ell_{I_\perp}^{\tau\perp}\rangle$ $(I = 1, 2)$; the former is $|\ell_I\rangle$ projected

[6] The chemical potential of the right-handed neutrinos are Boltzmann-suppressed and irrelevant for the conditions of the equilibrium [26, 27, 44], for instance, given as $\mu_{\ell_\alpha} + \mu_{\tilde{H}_u} + \mu_{\tilde{g}} = 0$. See also appendix B.

[7] We use an estimation of the variance $\sqrt{\langle\phi^2\rangle} \sim (\hat{y}\hat{y}^\dagger)_{33}/(4\pi)M_{\rm pl} \sim 10^{11}\,{\rm GeV} \times \left[(\hat{y}\hat{y}^\dagger)_{33}/10^{-6}\right]$, where $M_{\rm pl}$ is the reduced Planck mass.

[8] We consider no fine-tuning in K_I^α. Namely, we consider $K_I^\alpha \gg 1$.

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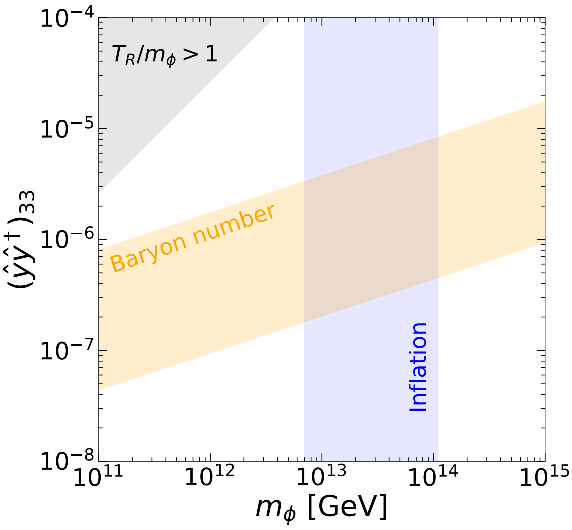

Figure 1. Parameter region which is consistent with the observed baryon asymmetry (orange, "Baryon number"). Here we take $ad = 0.01$ and 0.8 and use 1σ region of the observed baryon number given in eq. (1.1). Shaded region on the top-left corner indicate $T_R/m_\phi > 1$ (gray). We also plot the inflaton mass region (3.23), which is preferred the superconformal subcritical hybrid inflation [9, 10] (blue, "Inflation").

on to a plane perpendicular to $|\ell_\tau\rangle$ and the latter is one which is orthogonal to $|\ell_I^{\tau_\perp}\rangle$ in that plane. Then $|\ell_I^{\tau_\perp}\rangle$ component is washed out due to $K_I^e + K_I^\mu \gg 1$. To summarize, using the decomposition

$$|\ell_\phi\rangle = C_{\phi\tau}\,|\ell_\tau\rangle + C_{\phi\tau^\perp}\,|\ell_\tau^\perp\rangle \;, \tag{3.19}$$

$$|\ell_\tau^\perp\rangle = C_{\tau^\perp 2}\,|\ell_2^{\tau_\perp}\rangle + C_{\tau^\perp 2^\perp}\,|\ell_{2_\perp}^{\tau_\perp}\rangle \;, \tag{3.20}$$

$$|\ell_{2_\perp}^{\tau_\perp}\rangle = C_{2^\perp 1}\,|\ell_1^{\tau_\perp}\rangle + C_{2^\perp 1^\perp}\,|\ell_{1_\perp}^{\tau_\perp}\rangle \;, \tag{3.21}$$

a fraction $|C_{\phi\tau^\perp}C_{\tau^\perp 2^\perp}C_{2^\perp 1^\perp}|^2$ of the $B-L$ produced by the inflaton decay survives, i.e., $\mu_{Q\Delta_0} = |C_{\phi\tau^\perp}C_{\tau^\perp 2^\perp}C_{2^\perp 1^\perp}|^2 \mu_{L^{\rm dec}}$ [27]. $\mu_{Q\Delta_0}$ depends on the details of the model parameters, such as y_{ij}, M_{ij}, and λ_i.

If we consider $M_{1,2} \sim 10^6$–10^9 GeV, then all charged leptons are distinguished. Then all $B - L_\alpha$ components are washed out and no sufficient $B-L$ is obtained to explain the observed baryon number.

Finally at the stage *d)*, the $B-L$ is converted to the baryon number as $Y_B = (10/31)Y_{B-L}$. Combining eqs. (3.1), (3.10), (3.11), and (3.18), we get

$$Y_B \simeq 1.2\times 10^{-11}\left(\frac{ad}{0.01}\right)\left(\frac{m_\phi}{10^{13}\,\mathrm{GeV}}\right)^{-1/2}\left(\frac{(\hat{y}\hat{y}^\dagger)_{33}}{10^{-6}}\right)^{3/2} . \tag{3.22}$$

Here we take into account the spectator effects and the wash-out effect by introducing a coefficient d, which range from about 0.04 to 0.8. The result is plotted on $((\hat{y}\hat{y}^\dagger)_{33}, m_\phi)$ plane in figure 1. Here we also indicate the region where inflation induced by $\tilde{N}_3'$ predicts the spectral index and the tensor-to-scalar ratio that are consistent with the Planck observation

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based on refs. [9, 10] (see appendix C for details):

$$0.7 \times 10^{13}\,\mathrm{GeV} < m_\phi < 11 \times 10^{13}\,\mathrm{GeV}\,. \tag{3.23}$$

Therefore, the observed baryon number is obtained after the successful inflation in the region $(\hat{y}\hat{y}^\dagger)_{33} \sim 10^{-7}$–$10^{-6}$ and $m_\phi \sim 10^{13}\,\mathrm{GeV}$. This value of the neutrino Yukawa coupling is desirable for the reason discussed below. Let us say that all y_{ij} have the same order. Then we obtain $(\hat{y}\hat{y}^\dagger) \sim m_i\Lambda/\langle H_u\rangle^2 \sim 10^{-6}(\Lambda/10^{10}\,\mathrm{GeV})$ from the observed neutrino masses. This is consistent with the assumption (3.16). In addition, the baryon number is predicted to behave as $Y_B \propto \Lambda^{3/2}$, which means that $\Lambda \sim 10^{10}\,\mathrm{GeV}$, i.e. $M_{1,2} \sim 10^{10}\,\mathrm{GeV}$, is required to get the observed number. Therefore, in a case where all y_{ij} are the same order, the observed baryon asymetry can be obtained in the setup of this hybrid inflation model that is preferred by both the Planck observation and the neutrino masses.

4 Conclusion

We consider a model of supersymmetric hybrid inflation and study the reheating and baryogenesis after inflation. The model consists of three right-handed neutrinos N_i with the Majorana masses and two fields $S_\pm$ that are charged under a gauged U(1). A scalar component of a linear combination of the N_i plays the role of the inflaton, while the S_+ is the waterfall field. We focus on a case where the inflation lasts below the critical point value, called subcritical hybrid inflation. The inflaton mass should be $[0.7, 11] \times 10^{13}\,\mathrm{GeV}$ from the observations of the scalar spectral index and tensor-to-scalar ratio by the Planck collaboration. In addition, the scale Λ of the Majorana masses needs to be smaller than $\mathcal{O}(10^{10})\,\mathrm{GeV}$ in order not to disturb the inflationary dynamics. Therefore there is a hierarchy between the inflaton mass and the Majorana mass scale. In addition, the VEV of the waterfall field results in Dirac mass terms for the N_i and S_-.

After inflation, the inflaton decays to reheat the universe and at the same time several conserved quantities, including $B/3 - L_\alpha$, are provided. The conserved charges are broken and transported under the equilibrium conditions as the temperature drops down. Furthermore $B/3 - L_\alpha$ suffers from the wash-out due to the lighter right-handed (s)neutrinos. The wash-out is inevitably strong because of the observed neutrino masses and the special structure, i.e., Dirac and Majorana type, of the mass matrix of the right-handed neutrinos. In spite of the strong wash-out, a part of $B/3 - L_\alpha$ can survive for $\Lambda \sim 10^9$–$10^{10}\,\mathrm{GeV}$. Below that scale, though there are conserved quantities, such as $\mathcal{R}$, they have negligible contributions to the baryon number. Consequently, we found a successful baryogenesis for $10^9\,\mathrm{GeV} \lesssim \Lambda \lesssim 10^{10}\,\mathrm{GeV}$.

Acknowledgments

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A The interactions and charges

We construct a set of U(1) charges in the MSSM based on the technique given in ref. [28]. In our study we ignore the μ term for Higgses and the masses of gauginos by assuming μ and the supersymmetry breaking scale smaller than $\mathcal{O}(10^9)$ GeV.

In the MSSM, we have the following relevant fields [43–45]:

$$f = (e, \mu, \tau, \ell_e, \ell_\mu, \ell_\tau, u, c, t, d, s, b, Q_1, Q_2, Q_3, \tilde{H}_u, \tilde{H}_d, \tilde{g})\,, \tag{A.1}$$

$$b = (\tilde{e}, \tilde{\mu}, \tilde{\tau}, \tilde{\ell}_e, \tilde{\ell}_\mu, \tilde{\ell}_\tau, \tilde{u}, \tilde{c}, \tilde{t}, \tilde{d}, \tilde{s}, \tilde{b}, \tilde{Q}_1, \tilde{Q}_2, \tilde{Q}_3, H_u, H_d, g)\,, \tag{A.2}$$

where f are fermions and b indicates their bosonic partners. e, μ, τ, u, c, t, d, s, b are right-handed fields, and the rest are left-handed fields. The gauge interactions are equilibrium and all gauginos have the same chemical potential, denoted as $\tilde{g}$. The number density asymmetries are given by their chemical potentials μ_i $(i = f, b)$ as

$$n_i - \bar{n}_i = g_i \mu_i T^2/6\,, \tag{A.3}$$

where g_i are the multiplicities defined by

$$g_f = (1, 1, 1, 2, 2, 2, 3, 3, 3, 3, 3, 3, 6, 6, 6, 2, 2, 12)\,, \tag{A.4}$$

$$g_b = (2, 2, 2, 4, 4, 4, 6, 6, 6, 6, 6, 6, 12, 12, 12, 4, 4, 0)\,. \tag{A.5}$$

Due to the gauge interactions, the chemical potentials of bosonic partners are given by $\mu_b = \mu_f + \mu_{\tilde{g}}$ and $\mu_f - \mu_{\tilde{g}}$ for left-handed and right-handed fields, respectively. Following refs. [28, 44],[9] we introduce the interaction vectors:

$$y^t = (0, 0, 0, 0, 0, 0, 0, 0, -1, 0, 0, 0, 0, 0, 1, 1, 0, 1)\,, \tag{A.6}$$

$$y^{\rm SS} = (0, 0, 0, 0, 0, 0, -1, -1, -1, -1, -1, -1, 2, 2, 2, 0, 0, 6)\,, \tag{A.7}$$

$$y^{\rm WS} = (0, 0, 0, 1, 1, 1, 0, 0, 0, 0, 0, 0, 3, 3, 3, 1, 1, 4)\,, \tag{A.8}$$

$$y^b = (0, 0, 0, 0, 0, 0, 0, 0, 0, 0, 0, -1, 0, 0, 1, 0, 1, 1)\,, \tag{A.9}$$

$$y^\tau = (0, 0, -1, 0, 0, 1, 0, 0, 0, 0, 0, 0, 0, 0, 0, 0, 1, 1)\,, \tag{A.10}$$

$$y^{Q_{23}} = (0, 0, 0, 0, 0, 0, 0, 0, 0, 0, 0, 0, 0, -1, 1, 0, 0, 0)\,, \tag{A.11}$$

$$y^c = (0, 0, 0, 0, 0, 0, 0, -1, 0, 0, 0, 0, 0, 1, 0, 1, 0, 1)\,, \tag{A.12}$$

$$y^\mu = (0, -1, 0, 0, 1, 0, 0, 0, 0, 0, 0, 0, 0, 0, 0, 0, 1, 1)\,, \tag{A.13}$$

$$y^s = (0, 0, 0, 0, 0, 0, 0, 0, 0, 0, -1, 0, 0, 1, 0, 0, 1, 1)\,, \tag{A.14}$$

$$y^{Q_{12}} = (0, 0, 0, 0, 0, 0, 0, 0, 0, 0, 0, 0, -1, 1, 0, 0, 0, 0)\,, \tag{A.15}$$

$$y^d = (0, 0, 0, 0, 0, 0, 0, 0, 0, -1, 0, 0, 1, 0, 0, 0, 1, 1)\,, \tag{A.16}$$

$$y^u = (0, 0, 0, 0, 0, 0, -1, 0, 0, 0, 0, 0, 1, 0, 0, 1, 0, 1)\,, \tag{A.17}$$

$$y^e = (-1, 0, 0, 1, 0, 0, 0, 0, 0, 0, 0, 0, 0, 0, 0, 0, 1, 1)\,. \tag{A.18}$$

'SS' and 'WS' are strong and weak sphaleron processes, respectively, and the others are from the Yukawa interactions.[10] Using the interaction vectors, the equilibrium condition is

[9] See also ref. [46].

[10] $y^{Q_{23}}$ and $y^{Q_{12}}$ can be replaced by $y^{sb} = (0, 0, 0, 0, 0, 0, 0, 0, 0, 0, -1, 0, 0, 0, 1, 0, 1, 1)$ and $y^{ds} = (0, 0, 0, 0, 0, 0, 0, 0, 0, -1, 0, 0, 0, 1, 0, 0, 1, 1)$, respectively.

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given by $\mu_f \cdot y^{\rm int} = 0$ ($int = t, \mathrm{SS}, \mathrm{WS}, \cdots$). For example, $\mu_f \cdot y^t = -\mu_t + \mu_{Q_3} + \mu_{\tilde{H}_u} + \mu_{\tilde{g}} = 0$ etc. To describe the transportation of the chemical potentials, we introduce a set of charges for fermions f based on refs. [28, 44]. There are five charges which are anomaly free and conserved under the interactions listed above:

$$q_f^Y = (-1,-1,-1,-1/2,-1/2,-1/2,2/3,2/3,2/3,-1/3,-1/3,-1/3,1/6,1/6,1/6,1/2,-1/2,0)\,, \tag{A.19}$$
$$q_f^{\Delta_e} = (-1,0,0,-1,0,0,1/9,1/9,1/9,1/9,1/9,1/9,1/9,1/9,1/9,0,0,0)\,, \tag{A.20}$$
$$q_f^{\Delta_\mu} = (0,-1,0,0,-1,0,1/9,1/9,1/9,1/9,1/9,1/9,1/9,1/9,1/9,0,0,0)\,, \tag{A.21}$$
$$q_f^{\Delta_\tau} = (0,0,-1,0,0,-1,1/9,1/9,1/9,1/9,1/9,1/9,1/9,1/9,1/9,0,0,0)\,, \tag{A.22}$$
$$q_f^{\mathcal{R}} = (2,2,2,0,0,0,-4/9,-4/9,-4/9,14/9,14/9,14/9,-4/9,-4/9,-4/9,-1,1,1)\,, \tag{A.23}$$

corresponding to hypercharge, $\Delta_\alpha = B/3 - L_\alpha$ ($\alpha = e, \mu, \tau$) and $\mathcal{R}$ introduced in ref. [44].[11] The rest charges are broken at the onset of each interactions from (A.6) to (A.18):

$$q_f^t = (0,0,0,0,0,0,0,0,1,0,0,0,0,0,0,0,0,0)\,, \tag{A.24}$$
$$q_f^u = (0,0,0,0,0,0,1,0,0,0,0,0,0,0,0,0,0,0)\,, \tag{A.25}$$
$$q_f^B = (0,0,0,0,0,0,1/3,1/3,1/3,1/3,1/3,1/3,1/3,1/3,1/3,0,0,0)\,, \tag{A.26}$$
$$q_f^{R_{3b}^\chi} = (5,5,5,2,2,2,-3,-3,-3,2,2,5,-1,-1,-1,-3,2,1)\,, \tag{A.27}$$
$$q_f^\tau = (0,0,1,0,0,0,0,0,0,0,0,0,0,0,0,0,0,0)\,, \tag{A.28}$$
$$q_f^{B_1-B_2} = (0,0,0,0,0,0,1/3,-1/3,0,1/3,-1/3,0,1/3,-1/3,0,0,0,0)\,, \tag{A.29}$$
$$q_f^{R_{3c}^\chi} = (5,5,5,2,2,2,-3,0,-3,2,2,2,-1,-1,-1,-3,2,1)\,, \tag{A.30}$$
$$q_f^\mu = (0,1,0,0,0,0,0,0,0,0,0,0,0,0,0,0,0,0)\,, \tag{A.31}$$
$$q_f^{R_{3s}^\chi} = (5,5,5,2,2,2,-3,-3,-3,2,5,2,-1,-1,-1,-3,2,1)\,, \tag{A.32}$$
$$q_f^{2B_1-B_2-B_3} = (0,0,0,0,0,0,2/3,-1/3,-1/3,2/3,-1/3,-1/3,2/3,-1/3,-1/3,0,0,0)\,, \tag{A.33}$$
$$q_f^{R_{3d}^\chi} = (5,5,5,2,2,2,-3,-3,-3,5,2,2,-1,-1,-1,-3,2,1)\,, \tag{A.34}$$
$$q_f^{R_{3u}^\chi} = (5,5,5,2,2,2,0,-3,-3,2,2,2,-1,-1,-1,-3,2,1)\,, \tag{A.35}$$
$$q_f^e = (1,0,0,0,0,0,0,0,0,0,0,0,0,0,0,0,0,0)\,. \tag{A.36}$$

The charge assignments for bosons are given by $q_b^i = q_f^i + 1$ and $q_f^i - 1$ ($i = Y, \Delta_e, \Delta_\mu, \cdots$) for left-handed and right-handed fields, respectively, for $\mathcal{R}$ and R_{3q}^χ ($q = b, c, s, d, u$), and they are the same as fermions for the others. We have introduced a modified version of charge, denoted as R_{3q}^χ, based on the *chiral* R_3 charge in ref. [44]. With above definition and eqs. (A.4) and (A.5), the chemical potentials of the conserved charges are given by

$$\mu_{Q^i} = (q_f^i \circ g_f) \cdot \mu_f + (q_b^i \circ g_b) \cdot \mu_b\,, \tag{A.37}$$

where $\circ$ denotes the entrywise Hadamard product. We always impose $\mu_{Q^Y} = 0$. With the chemical potentials μ_{Q^i}, the asymmetry number density for the conserved charges are

[11] This $\mathcal{R}$ is different from the R-symmetry. As a reference, R-charges of the fermions are given as $q_f^R = (-1,-1,-1,-1,-1,-1,3,3,3,3,3,3,-1,-1,-1,3,-1,1)$ [44].

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given by

$$n^{Q^i} - \bar{n}^{Q^i} = \mu_{Q^i} T^2/6\,. \tag{A.38}$$

For example, the asymmetry of $B-L$ is given by $\sum_\alpha \mu_{Q^{\Delta_\alpha}} T^2/6$.

B $\quad B-L$ at the wash-out regime

With the interactions and the conserved charges introduced in the previous section, we can calculate the $B-L$ for a given set of the equilibrium conditions. Here we give several temperature regime from $10^{10}\,\mathrm{GeV}$ to $10^5\,\mathrm{GeV}$. Though we focus on the range $[10^7\,\mathrm{GeV}, 10^{10}\,\mathrm{GeV}]$, we give the result as a reference. In out study, we consider $\tan\beta$, the ratio of the VEV of the up-type and down-type Higgs, is ~ 1 and adopt the equilibrium temperatures of the relevant interactions given in refs. [29, 47]. To check our calculation, we compute A^ℓ, $C^{\tilde{g}}$, $C^{\tilde{H}_u}$, and $C^{\tilde{H}_d}$ in ref. [44] and obtain consistent results from eqs. (2.40)–(2.45) in the literature, except for eq. (2.41).[12]

(i) $T \sim 10^9$–10^{10} GeV. t, b, c, τ Yukawa interactions and strong, weak sphaleron processes are in equilibrium. In this case, τ (including ℓ_τ) is distinguished. On the other hand, a linear combination of ℓ_e and ℓ_μ, which are tentatively denoted as ℓ'_e and ℓ'_μ, are disentangled if the interaction with $\hat{N}_1$ and $\hat{N}_2$ are in equilibrium, which will be discussed later.

It is straightforward to compute the chemical potentials of ℓ'_e, ℓ'_μ, ℓ_τ, $\tilde{H}_u$, and $\tilde{g}$ in terms of μ_{Q^i} $(i = \Delta_{e'}, \Delta_{\mu'}, \Delta_\tau, \mathcal{R}, e', \mu', R^\chi_{3u}, R^\chi_{3d}, 2B_1 - B_2 - B_3, R^\chi_{3s}, B_1 - B_2)$. The coefficients of μ_{Q^i} are given by

$$\mu_{\ell'_e}\colon\ -\tfrac{432337}{3164482}, \tfrac{142615}{4746723}, \tfrac{100513}{4746723}, -\tfrac{132149}{9493446}, -\tfrac{1329629}{9493446}, \tfrac{42102}{1582241}, \tfrac{4943}{1582241}, \tfrac{297}{93073}, -\tfrac{10045}{1582241}, \tfrac{10204}{4746723}, \tfrac{7653}{1582241} \tag{B.1}$$

$$\mu_{\ell'_\mu}\colon\ \tfrac{142615}{4746723}, -\tfrac{432337}{3164482}, \tfrac{100513}{4746723}, -\tfrac{132149}{9493446}, \tfrac{42102}{1582241}, -\tfrac{1329629}{9493446}, \tfrac{4943}{1582241}, \tfrac{297}{93073}, -\tfrac{10045}{1582241}, \tfrac{10204}{4746723}, \tfrac{7653}{1582241} \tag{B.2}$$

$$\mu_{\ell_\tau}\colon\ \tfrac{167309}{9493446}, \tfrac{167309}{9493446}, -\tfrac{166914}{1582241}, -\tfrac{273605}{9493446}, \tfrac{341897}{9493446}, \tfrac{341897}{9493446}, \tfrac{11213}{3164482}, \tfrac{6013}{558438}, -\tfrac{56717}{3164482}, \tfrac{15168}{1582241}, \tfrac{34128}{1582241} \tag{B.3}$$

$$\mu_{\tilde{H}_u}\colon\ -\tfrac{94921}{3164482}, -\tfrac{94921}{3164482}, -\tfrac{52238}{1582241}, \tfrac{46760}{1582241}, \tfrac{28665}{3164482}, \tfrac{28665}{3164482}, -\tfrac{203327}{9493446}, -\tfrac{13679}{1675314}, \tfrac{72645}{3164482}, -\tfrac{14608}{14240169}, -\tfrac{3652}{1582241} \tag{B.4}$$

$$\mu_{\tilde{g}}\colon\ -\tfrac{17141}{1582241}, -\tfrac{17141}{1582241}, -\tfrac{17983}{1582241}, \tfrac{45461}{3164482}, \tfrac{2526}{1582241}, \tfrac{2526}{1582241}, \tfrac{1424}{1582241}, \tfrac{332}{279219}, -\tfrac{3534}{1582241}, \tfrac{4220}{4746723}, \tfrac{3165}{1582241}\,. \tag{B.5}$$

Now we take into account the strong wash-out effect. The chemical potentials after the wash-out are given by $\mu_{\ell_1} + \mu_{\tilde{H}_u} + \mu_{\tilde{g}} = 0$ and $\mu_{\ell_2} + \mu_{\tilde{H}_u} + \mu_{\tilde{g}} = 0$. Here it should be noted that a charged lepton $\ell^{\tau_\perp}_{1_\perp}$ in eq. (3.21) does not couple to $\hat{N}_1$ and $\hat{N}_2$ and the corresponding charge is conserved and the others are broken. However, it is not trivial to extract such a state from the equilibrium equations and solve them analytically. Following ref. [26], we approximately estimate the wash-out effect by solving $\mu_{\ell_\tau} + \mu_{\tilde{H}_u} + \mu_{\tilde{g}} = 0$ and $\mu_{\ell'_\mu} + \mu_{\tilde{H}_u} + \mu_{\tilde{g}} = 0$ and identify ℓ'_e as $\ell^{\tau_\perp}_{1_\perp}$. A crucial point is that the neutrino Yukawa interactions do not break the other symmetries. Namely, $\Delta_0, \mathcal{R}, e', \mu', R^\chi_{3u}, R^\chi_{3d}, 2B_1 - B_2 - B_3, R^\chi_{3s}$, and $B_1 - B_2$ are unbroken. Here we rewrite $\Delta_{e'}$ as Δ_0. Therefore, in the basis μ_{Q^i} $(i = \Delta_0, \mathcal{R}, e', \mu', R^\chi_{3u}, R^\chi_{3d}, 2B_1 - B_2 - B_3, R^\chi_{3s}, B_1 - B_2)$ the coefficients to give chemical potential for $B-L$ are given by

$$\mu_{Q^{B-L}}\ \colon\ \tfrac{120251}{148420}, \tfrac{3491}{14842}, \tfrac{67889}{148420}, -\tfrac{26181}{74210}, -\tfrac{27463}{148420}, \tfrac{1801}{445260}, \tfrac{12831}{148420}, \tfrac{7316}{111315}, \tfrac{5487}{37105}\,. \tag{B.6}$$

[12] We thank Chee Sheng Fong for confirming this point.

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(ii) $T \sim 10^6$–10^9 GeV. In addition to the previous case, μ, s, Q_{23} Yukawa interactions are in equilibrium. Then all charged leptons are disentangled. The results are $(i = \Delta_e, \Delta_\mu, \Delta_\tau, \mathcal{R}, e, R^\chi_{3u}, R^\chi_{3d}, 2B_1 - B_2 - B_3)$. The coefficients of μ_{Q^i} are given by

$$\mu_{\ell_e}: \quad -\tfrac{49715}{369014}, \tfrac{13217}{553521}, \tfrac{13217}{553521}, -\tfrac{11719}{1107042}, -\tfrac{157723}{1107042}, \tfrac{1759}{553521}, \tfrac{1015}{553521}, -\tfrac{1387}{369014} \tag{B.7}$$
$$\mu_{\ell_\mu}: \quad \tfrac{45005}{2214084}, -\tfrac{165299}{1660563}, \tfrac{19208}{1660563}, -\tfrac{19111}{1107042}, \tfrac{58183}{2214084}, \tfrac{9695}{2214084}, \tfrac{5571}{738028}, -\tfrac{3301}{369014} \tag{B.8}$$
$$\mu_{\ell_\tau}: \quad \tfrac{45005}{2214084}, \tfrac{19208}{1660563}, -\tfrac{165299}{1660563}, -\tfrac{19111}{1107042}, \tfrac{58183}{2214084}, \tfrac{9695}{2214084}, \tfrac{5571}{738028}, -\tfrac{3301}{369014} \tag{B.9}$$
$$\mu_{\tilde{H}_u}: \quad -\tfrac{21687}{738028}, -\tfrac{6033}{184507}, -\tfrac{6033}{184507}, \tfrac{5320}{184507}, \tfrac{7335}{738028}, -\tfrac{143473}{6642252}, -\tfrac{54887}{6642252}, \tfrac{8265}{369014} \tag{B.10}$$
$$\mu_{\tilde{g}}: \quad -\tfrac{1975}{184507}, -\tfrac{2019}{184507}, -\tfrac{2019}{184507}, \tfrac{5671}{369014}, \tfrac{132}{184507}, \tfrac{548}{553521}, \tfrac{526}{553521}, -\tfrac{537}{369014}\,. \tag{B.11}$$

As the previous case, we solve $\mu_{\ell_\alpha} + \mu_{\tilde{H}_u} + \mu_{\tilde{g}} = 0$ $(\alpha = e, \mu, \tau)$ to take into account the wash-out effect. In this case all Δ_α are broken due to the neutrino Yukawa interactions. Consequently, the chemical potential for $B - L$ is given by the chemical potentials of the unbroken charges μ_{Q^i} $(i = \mathcal{R}, e, R^\chi_{3u}, R^\chi_{3d}, 2B_1 - B_2 - B_3)$,

$$\mu_{Q^{B-L}}: \quad \tfrac{85794}{198313}, -\tfrac{42771}{198313}, -\tfrac{49149}{198313}, -\tfrac{4455}{198313}, \tfrac{40203}{198313}\,. \tag{B.12}$$

(iii) $T \sim 10^5$–10^6 GeV. Under this temperature only e Yukawa interaction is out of equilibrium. The coefficients of μ_{Q^i} $(i = \Delta_e, \Delta_\mu, \Delta_\tau, \mathcal{R}, e)$ are given by

$$\mu_{\ell_e}: \quad -\tfrac{931}{6786}, \tfrac{211}{9657}, \tfrac{211}{9657}, -\tfrac{1}{174}, -\tfrac{415}{2886} \tag{B.13}$$
$$\mu_{\ell_\mu}: \quad \tfrac{113}{6786}, -\tfrac{326}{3219}, \tfrac{95}{9657}, -\tfrac{1}{174}, \tfrac{59}{2886} \tag{B.14}$$
$$\mu_{\ell_\tau}: \quad \tfrac{113}{6786}, \tfrac{95}{9657}, -\tfrac{326}{3219}, -\tfrac{1}{174}, \tfrac{59}{2886} \tag{B.15}$$
$$\mu_{\tilde{H}_u}: \quad -\tfrac{1}{78}, -\tfrac{2}{111}, -\tfrac{2}{111}, 0, \tfrac{15}{962} \tag{B.16}$$
$$\mu_{\tilde{g}}: \quad -\tfrac{1}{87}, -\tfrac{1}{87}, -\tfrac{1}{87}, \tfrac{1}{58}\,. \tag{B.17}$$

By solving $\mu_{\ell_\alpha} + \mu_{\tilde{H}_u} + \mu_{\tilde{g}} = 0$ $(\alpha = e, \mu, \tau)$, the chemical potential for $B - L$ is given by

$$\mu_{Q^{B-L}}: \quad \tfrac{213}{1012}, -\tfrac{261}{1012} \tag{B.18}$$

in the basis μ_{Q^i} $(i = \mathcal{R}, e)$.

C Inflaton mass

We use the latest Planck data [1, 48]

$$n_s = 0.9649 \pm 0.0042 \, (68\% \, \text{C.L.}) \,, \tag{C.1}$$
$$r < 0.10 \, (95\% \, \text{C.L.}) \,, \tag{C.2}$$
$$A_s = 2.100 \pm 0.030 \times 10^{-9} \, (68\% \, \text{C.L.}) \,, \tag{C.3}$$

where n_s, r, and A_s are the spectral index of the scalar mode, tensor-to-scalar ratio, and the amplitude of the scalar mode. Based on the analysis given in ref. [10], we evaluate the inflaton mass. The result is shown in figure 2, where the minimum and maximum values

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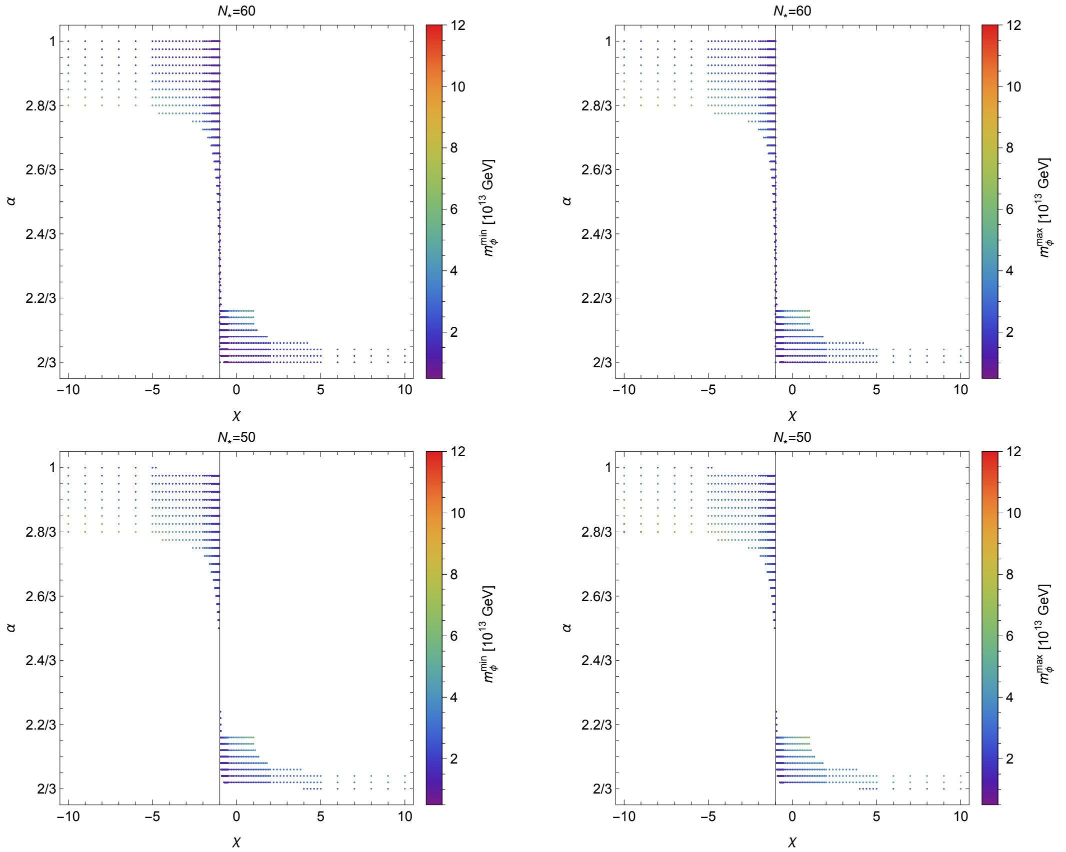

Figure 2. Minimum (left) and maximum (right) values of the inflaton mass for a given set of parameters, α and χ, that characterize the symmetry of the Kähler potential of the model [10]. We take the number of e-folds as 60 (top) and 50 (bottom).

of the inflaton mass is indicated for a given set of the model parameters. We found the inflaton mass range that is consistent with the Planck result is

$$0.7 \times 10^{13}\,\mathrm{GeV} < m_\phi < 10 \times 10^{13}\,\mathrm{GeV}\,, \tag{C.4}$$

for 60 e-folds and

$$1.2 \times 10^{13}\,\mathrm{GeV} < m_\phi < 11 \times 10^{13}\,\mathrm{GeV}\,, \tag{C.5}$$

for 50 e-folds.

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