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Leptogenesis after superconformal subcritical hybrid inflation

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ABSTRACT: We consider an extended version of superconformal subcritical hybrid inflation model by introducing three right-handed neutrinos that have the Majorana mass terms. In the model one of the right-handed sneutrinos plays a role of the inflaton field, and it decays to reheat the universe after inflation. The vacuum expectation value for the waterfall field gives an unconventional pattern of the light neutrino mass matrix, and the neutrino Yukawa couplings that determine the reheating temperature are not constrained by the neutrino oscillation data. Consequently thermal leptogenesis or sneutrino leptogenesis is realized.

KEYWORDS: Cosmology of Theories beyond the SM, CP violation, Neutrino Physics, Supergravity Models

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1 Introduction

Inflation paradigm is strongly supported by the observations of the cosmic microwave background radiation (CMB). Slow-roll scalar field in the early universe is a promising candidate for inflation, and many types of inflation models have been proposed so far. In a theoretical point of view, it would be tempting to ask what the underlying physics or symmetry of the inflaton field is. Supersymmetry (SUSY) might be one of the answers. It protects the flatness of the inflaton direction, which is suitable for inflation.

Recently supersymmetric D-term hybrid inflation has been revisited in various point of view. Under shift symmetric Kähler potential [1], subcritical hybrid inflation was found, where inflation continues for subcritical point value of the inflaton field [2, 3]. On the other hand, it was shown in refs. [4, 5] that Starobinsky model [6] emerges in the framework of superconformal supergravity [7–10]. It turned out in the following study that this framework has another new regime of inflation. It was shown that a general class of superconformal α -attractor model [11, 12] appears in the subcritical regime of inflation, which we call superconformal subcritical hybrid inflation [13]. In addition, the energy scale of inflation should coincide with the grand unification scale to be consistent with the Planck observation data, which is the feature found in the subcritical hybrid inflation [2, 3].

Namely, the superconformal subcritical hybrid inflation has both features of the superconformal α -attractor models and the subcritical hybrid inflation. The shift symmetry and superconformality are crucial for them.

In this paper we will study the thermal history after the end of superconformal subcritical hybrid inflation. (See refs. [14, 15] that study the phenomenology of Pati-Salam version of subcritical hybrid inflation. Recently ref. [16] comprehensively studies the D-term hybrid inflation, including reheating, leptogenesis, and the SUSY breaking mechanism.) For the purpose, we introduce three right-handed neutrinos that interact with the minimal supersymmetric standard model (MSSM) sector. In fermionic sector, the mass matrix for the light neutrinos is given by the seesaw mechanism [17–22], but it has an unconventional structure. In bosonic sector, on the other hand, it will be shown that one of sneutrinos can play a role of the inflaton field. In addition, baryon asymmetry that is sufficient amount to explain the observed value is generated via leptogenesis [23].

This paper is organized as follows. In the next section the model that we consider is described. Then section 3 shows the conditions required for the superconformal subcritical hybrid inflation in this model. Mass matrices of the heavy and light (s)neutrinos, including parametrization of the neutrino Yukawa couplings, are given in section 4, then we discuss the reheating and leptogenesis after inflation in section 5. Section 6 is dedicated to conclusion.

2 The model

We consider a model described by the superpotential

$$W = W_{\text{MSSM}} + W_{\text{neu}} , \quad (2.1)$$

where W_{MSSM} is the superpotential of the MSSM sector and

$$W_{\text{neu}} = \frac{1}{2} M_{ij} N_i^c N_j^c + y_{\nu ij} N_i^c L_j H_u + \lambda_i N_i^c S_+ S_- . \quad (2.2)$$

Here N_i^c , $L_i = (\nu_{Li}, l_{Li})^T$, and $H_u = (H_u^+, H_u^0)^T$ are the chiral superfields of the right-handed neutrinos, left-handed leptons, and up-type Higgs, respectively, and $L_j H_u = \nu_{Lj} H_u^0 - l_{Lj} H_u^+$. $S_{\pm}$ are the local U(1) fields with charge $\pm q$ ($q > 0$), one of which plays a role of the waterfall field.¹ Indices are summed over $i, j = 1-3$. (We will use the same contraction in the following discussion unless otherwise mentioned.) M_{ij} terms are explicit superconformal breaking terms that are added by phenomenological purpose. While we do not argue their origin, $M_{ij} \ll 1$ are expected. The Kähler potential, on the other hand, is given by

$$\mathcal{K} = -3 \log \Omega^{-2} , \quad (2.3)$$

¹We write the superpartners with tilde for the MSSM fields and right-handed neutrinos. For $S_{\pm}$, the same symbols are used for scalar fields while fermionic parts are expressed with tilde. In the current and the next sections, we adopt the unit in which the reduced Planck mass $M_{pl} \simeq 2.4 \times 10^{18}$ GeV is taken to be unity unless otherwise mentioned.

where

$$\Omega^{-2} = 1 - \frac{1}{3} \left(\sum_{\text{MSSM}} |I_{\text{MSSM}}|^2 + \sum_i |N_i^c|^2 + |S_+|^2 + |S_-|^2 \right) - \sum_i \frac{\chi_i}{6} (N_i^{c2} + \bar{N}_i^{c2}) . \quad (2.4)$$

Here I_{MSSM} are chiral superfields in the MSSM sector. The last term is superconformal breaking term that is considered in refs. [4, 10]. With the superpotential and Kähler potential, the scalar potential is given by

$$V_{\text{tot}} = V_F + V_D, \quad (2.5)$$

where V_F and V_D are F- and D-terms, respectively, and given by [4]

$$V_F = \Omega^4 \left[\delta^{\alpha\bar{\alpha}} W_{\alpha} W_{\bar{\alpha}} + \frac{1}{\Delta} |\delta^{\alpha\bar{\alpha}} W_{\alpha} \Phi_{\bar{\alpha}} - 3W|^2 \right], \quad (2.6)$$

$$V_D = \frac{1}{2} g^2 [q\Omega^2 (|S_+|^2 - |S_-|^2) - \xi]^2 . \quad (2.7)$$

Here $\Phi \equiv -3\Omega^{-2}$ and $\Delta \equiv \Phi - \delta^{\alpha\bar{\alpha}} \Phi_{\alpha} \Phi_{\bar{\alpha}}$ have been additionally introduced. Subscript in W and Φ stands for the field derivative, e.g., $W_{\alpha} \equiv \partial W / \partial z^{\alpha}$ where z^{α} is a chiral superfield. In the D-term, we have introduced the Fayet-Iliopoulos (FI) term $\xi (> 0)$ associated with the U(1).² Due to the FI term, S_+ has a vacuum expectation value (VEV) at the global minimum, which is obtained as $\langle S_+ \rangle = \sqrt{\xi/q(1+\tilde{\xi})}$ with $\tilde{\xi} \equiv \xi/3q$.

As in ref. [13], we take $\chi_i \leq 0$ without loss of generality. In the present model, we impose the following condition:

$$\chi_3 \simeq -1, \quad \chi_1, \chi_2 \simeq 0. \quad (2.8)$$

This distinguishes N_3^c from N_1^c and N_2^c . $\phi \equiv \sqrt{2} \text{Re } \tilde{N}_3^c$ has an approximate shift symmetry that is explicitly broken by $\lambda_3 (\ll 1)$. Then, ϕ is expected to be the inflaton as studied in ref. [13]. In the inflation model, $s \equiv \sqrt{2}|S_+|$ plays a role of the waterfall field. N_1^c and N_2^c , on the other hand, have no such symmetry. Instead there is a freedom to choose any basis for N_1^c and N_2^c with a redefinition of M_{ij} and λ_i due to $\chi_{1,2} \simeq 0$.

For simple notation, hereafter we omit subscripts of χ_3 and λ_3 and introduce m_{ϕ} , which will be identified as the inflaton mass in section 3.2 (with another assumption in section 4.1),

$$\chi \equiv \chi_3, \quad \lambda \equiv \lambda_3, \quad (2.9)$$

$$m_{\phi} \equiv \lambda \langle S_+ \rangle. \quad (2.10)$$

It is sometimes convenient to use $\delta\chi$ ($0 < \delta\chi < 1$) defined by $\delta\chi/(1+\chi) = -qg^2\xi/3\lambda^2$. $0 < \delta\chi < 1$ guarantees $\Omega(\phi, s)^2 > 0$ and $\phi_{c,0}^2 > 0$, which will be defined in eqs. (3.2)

²The origin of the FI term in canonical superconformal supergravity model [10] is discussed in ref. [4]. See also ref. [16] for recent development.

and (3.4), respectively. Then inflation that is consistent with the observations of the CMB is realized in the parameter space [13],

$$\begin{aligned}\lambda &\simeq (0.5-1) \times 10^{-3}, & \xi^{1/2} &\simeq (3-1) \times 10^{16} \text{ GeV}, \\ m_\phi &\simeq (1-2) \times 10^{13} \text{ GeV},\end{aligned}\tag{2.11}$$

for $q = g = 1$, $\delta\chi = 0.9$ and the number of e -folds $N_e = 55-60$, which we take in the later numerical study.³ In the parameter space, e.g., $\chi \simeq -1.16$ for $\lambda = 10^{-3}$, $\xi^{1/2} = 10^{16} \text{ GeV}$, and $\delta\chi = 0.9$. The mass terms in the superpotential, however, have a possibility to alter the inflationary path. In the next section, we will derive the conditions in order not to affect the inflationary dynamics.

3 Inflation

We define several variables that are used in the following analysis. During inflation, the other fields except for the inflaton and waterfall field are irrelevant. Thus it is convenient to define following potentials,

$$V(\phi, s) \equiv V_{\text{tot}}|_{\sqrt{2}\text{Re } \tilde{N}_3^c = \phi, \sqrt{2}|S_+| = s, \text{ the others}=0},\tag{3.1}$$

$$\Omega(\phi, s) \equiv \Omega|_{\sqrt{2}\text{Re } \tilde{N}_3^c = \phi, \sqrt{2}|S_+| = s, \text{ the others}=0}.\tag{3.2}$$

Then the critical point value ϕ_c is defined as a field value below which the waterfall field becomes tachyonic. It receives $\mathcal{O}(M_{ij}^2)$ corrections as

$$\phi_c^2 = \phi_{c,0}^2 \left[1 - \frac{2\Delta M^2(\phi_{c,0})}{3\lambda^2} \right] + \mathcal{O}(M_{i3}^4),\tag{3.3}$$

where

$$\phi_{c,0}^2 = \frac{6qg^2\xi}{3\lambda^2 + (1+\chi)qg^2\xi},\tag{3.4}$$

$$\Delta M^2(\phi) = \sum_{i=1}^3 |M_{i3}|^2 - \frac{(1-2\chi)^2}{24} \frac{\phi^2}{1 + \phi^2\chi(1+\chi)/6} |M_{33}|^2.\tag{3.5}$$

For example, $\phi_{c,0} \simeq 18$ or in terms of canonically-normalized field $\hat{\phi}_{c,0} \simeq 11$ defined in eq. (3.12) for $\lambda = 10^{-3}$, $\xi^{1/2} = 10^{16} \text{ GeV}$ and $\delta\chi = 0.9$. This perturbative expansion is valid when⁴

$$|M_{13}|^2 + |M_{23}|^2 \ll 3\lambda^2/2 \sim (10^{15} \text{ GeV})^2,\tag{3.6}$$

$$|M_{33}|^2 \ll 2\lambda^4/qg^2\xi \sim (10^{14} \text{ GeV})^2.\tag{3.7}$$

It will be checked in this section that the above conditions are satisfied in this inflation model.

³We will estimate the number of e -folds in section 5.1 to confirm this.

⁴We have checked that $\mathcal{O}(M_{i3}^4)$ term is irrelevant when eq. (3.7) is satisfied, thus we ignore it in the following discussion.

Finally it is useful to define

$$\Psi \equiv \frac{\Omega(\phi, 0)\phi}{\Omega(\phi_{c,0}, 0)\phi_{c,0}} = \frac{\Omega(\phi, 0)\phi}{\sqrt{2qg^2\xi/\lambda^2}}, \quad (3.8)$$

when potential is expressed in terms of canonically-normalized inflaton field.

3.1 Pre-critical regime

Let us begin with the regime where the inflaton is approaching down to the critical point value. Since the waterfall field is stabilized at the origin in this regime, the relevant Lagrangian is given as

$$\mathcal{L}_{\text{pre}} = \frac{f(\phi, 0)}{2}(\partial_\mu\phi)^2 - V_{\text{pre}}(\phi), \quad (3.9)$$

where⁵

$$f(\phi, s) = \Omega^2(\phi, s) \left[1 + \Omega^2(\phi, s) \frac{(1+\chi)^2}{6} \phi^2 \right], \quad (3.10)$$

$$V_{\text{pre}}(\phi) = \frac{1}{2}g^2\xi^2 \left[1 + 2\Psi^2 \frac{\Omega^2(\phi, 0)\Delta M^2(\phi)}{\lambda^2\xi/q} \right]. \quad (3.11)$$

It is seen that ΔM^2 term gives a gradient to the inflaton field, which should not invade the slow-roll conditions. To see the impact of ΔM^2 term, it is instructive to change dynamical variable ϕ to canonically-normalized field $\hat{\phi}$. Since $\chi \simeq -1$, it is a good approximation that $f(\phi, 0) \simeq \Omega^2(\phi, 0)$. Then $d\phi/d\hat{\phi} = f(\phi)^{-1/2} \simeq \Omega^{-1}(\phi, 0)$ can be solved easily to obtain,

$$\phi \simeq \beta^{-1/2} \sinh \beta^{1/2} \hat{\phi}, \quad (3.12)$$

$$\Psi \simeq \delta\chi^{-1/2} \tanh \beta^{1/2} \hat{\phi}, \quad (3.13)$$

$$\Omega^{-1}(\phi, 0) \simeq \cosh \beta^{1/2} \hat{\phi}, \quad (3.14)$$

where $\beta = -(1+\chi)/6 = \lambda^2\delta\chi/2qg^2\xi$. Then the potential in terms of $\hat{\phi}$ is given as

$$V_{\text{pre}} \simeq \frac{1}{2}g^2\xi^2 \left[1 + \frac{\tanh^2 \beta^{1/2} \hat{\phi}}{\cosh^2 \beta^{1/2} \hat{\phi}} \frac{2q}{\lambda^2\xi\delta\chi} \left\{ \sum_{i=1}^3 |M_{i3}|^2 - \frac{3}{8\beta} |M_{33}|^2 \tanh^2 \beta^{1/2} \hat{\phi} \right\} \right]. \quad (3.15)$$

On the other hand, it was shown in ref. [4] that there is one-loop corrections to the tree-level potential. In terms of the canonically-normalized field, it is given by

$$V_{1l} \simeq \frac{1}{2}g^2\xi^2 \times \frac{q^2g^2}{8\pi^2} \log \left[\delta\chi^{-1} \tanh^2 \beta^{1/2} \hat{\phi} \right]. \quad (3.16)$$

Therefore, in order not to affect the inflationary trajectory, it is sufficient that the terms proportional to $|M_{i3}|^2$ are subdominant compared to the one-loop potential. In the parameter space given in eq. (2.11), $\sqrt{\beta}\hat{\phi}_c \simeq \text{arcsinh}\sqrt{\beta}\phi_c \sim \mathcal{O}(1)$. Then, the conditions are

⁵It is noted that the term proportional to ΔM^2 in eq. (3.11) is equivalent to $\Omega^4(\phi, 0)\phi^2\Delta M^2(\phi)/2$.

given as

$$|M_{13}|^2 + |M_{23}|^2 \lesssim \frac{qg^2\delta\chi\lambda^2\xi}{16\pi^2} \lesssim (1 \times 10^{12} \text{ GeV})^2, \quad (3.17)$$

$$|M_{33}|^2 \lesssim \frac{\delta\chi^2\lambda^4}{12\pi^2} \lesssim (2 \times 10^{11} \text{ GeV})^2. \quad (3.18)$$

It is easy to check that the slow-roll conditions are satisfied under the constraints. Since the constraints are more stringent than eqs. (3.6) and (3.7), it has been confirmed that the perturbative expansion to obtain eq. (3.3) is valid.

3.2 Subcritical regime

In the previous subsection, we have seen that the slow-roll conditions are satisfied before reaching to the critical point value. After the inflaton field becomes subcritical point value, the tachyonic growth of the waterfall field occurs. It is expected that the inflation continues in the subcritical regime when $M_{i3} \rightarrow 0$. In this subsection, we will derive the conditions under which the inflaton and waterfall field dynamics are not affected with non-zero M_{i3} .

As seen in the previous section, the perturbative expression for ϕ_c is valid under the conditions given in eqs. (3.17) and (3.18). Then, the tachyonic growth of the waterfall field is not affected by the additional gradient in the inflaton direction due to $|M_{i3}|^2$ terms since $\phi_c \simeq \phi_{c,0}$. Consequently, the dynamics of the waterfall field around the critical point is the same as one discussed in ref. [13]. The tachyonic growth is qualitatively the same as the subcritical hybrid inflation [2, 3]. Namely, due to the tachyonic growth, the waterfall field relaxes to the local minimum value $s_{\min}$ just after a few Hubble-unit time. (See, for example, figure 1 of ref. [2].) In the present model, $s_{\min}$ is found to be

$$s_{\min}^2 = \frac{2\xi}{q(1+\tilde{\xi})} \frac{\Omega^{-2}(\phi, 0)}{1 + \tilde{\xi}\Psi^2/(1+\tilde{\xi})} \left[1 - \Psi^2 \left\{ 1 + \frac{2\Omega^2(\phi, 0)\Delta M^2(\phi)}{3\lambda^2} \right\} \right]. \quad (3.19)$$

Then the potential in the subcritical regime of the inflaton field is effectively given by $V(\phi, s_{\min})$ and the dynamics reduces to single field inflation that is described by the Lagrangian,

$$\mathcal{L} = \frac{f(\phi, s_{\min})}{2} (\partial_\mu \phi)^2 - V_{\text{inf}}(\phi), \quad (3.20)$$

where

$$V_{\text{inf}}(\phi) = g^2\xi^2 \frac{(1+\tilde{\xi})\Psi^2}{1+2\tilde{\xi}\Psi^2} \left[1 - \frac{\Psi^2}{2(1+\tilde{\xi})} + \frac{\{1+\tilde{\xi}\Psi^2/(1+\tilde{\xi})\}^2 \Omega^2(\phi, 0)\Delta M^2(\phi)}{(1+2\tilde{\xi}\Psi^2)/(1+\tilde{\xi}) \lambda^2\xi/q} \right] + \mathcal{O}(M_{i3}^4). \quad (3.21)$$

Recall that $\tilde{\xi}, \xi \ll 1$ in the parameters in eq. (2.11). Then it is clear that the additional term proportional to ΔM^2 is the same as in eq. (3.11). Note that $s_{\min}$ is negligible in $\Omega(\phi, s_{\min})$. Then, the canonically-normalized inflaton field is given by eqs. (3.12)–(3.14).

Consequently, V_{inf} is given by

$$V_{\text{inf}} \simeq g^2 \xi^2 \delta \chi^{-1} \tanh^2 \beta^{1/2} \hat{\phi} \left[1 - \frac{\delta \chi^{-1}}{2} \tanh^2 \beta^{1/2} \hat{\phi} + \frac{\cosh^{-2} \beta^{1/2} \hat{\phi}}{\lambda^2 \xi / q} \left\{ \sum_{i=1}^3 |M_{i3}|^2 - \frac{3}{8\beta} |M_{33}|^2 \tanh^2 \beta^{1/2} \hat{\phi} \right\} \right]. \quad (3.22)$$

Therefore, if eqs. (3.17) and (3.18) are satisfied, then the dynamics in the subcritical regime reduces to one in ref. [13]. In addition, it is clear that m_ϕ defined in eq. (2.10) is indeed the inflaton mass since $V_{\text{inf}} \simeq \frac{1}{2} m_\phi^2 \phi^2$ around $\phi (\simeq \hat{\phi}) \sim 0$.

To summarize the present and previous subsections, the inflaton-waterfall field dynamics is unchanged when

$$|M_{i3}|^2 = |M_{3i}|^2 \lesssim (2 \times 10^{11} \text{ GeV})^2, \quad (3.23)$$

for $i = 1-3$ are satisfied.

3.3 Stability of inflationary trajectory

It was pointed out in ref. [24] that $\tilde{L}_i H_u$ may become tachyonic in sneutrino inflation. In order to find out the stability condition, let us derive the mass matrix in $\tilde{L}_i$ and H_u basis. From V_{tot} , it is obtained by

$$\begin{aligned} \mathcal{L}_{\text{mass}}^{\tilde{L}-H_u} = & -\frac{\Omega(\phi, s)^4 \phi^2}{2} |y_{\nu 3i} \tilde{L}_i|^2 - \frac{\Omega(\phi, s)^4 \phi^2}{2} |y_{\nu 3i}|^2 |H_u|^2 \\ & + \left[\frac{M_{33}^* \Omega(\phi, s)^6 \phi^3}{4\sqrt{2}} y_{\nu 3i} \tilde{L}_i H_u + \text{h.c.} \right]. \end{aligned} \quad (3.24)$$

On the other hand, the kinetic terms of $\tilde{L}_i$ and H_u are given by

$$\mathcal{L}_{\text{kin}}^{\tilde{L}-H_u} = \Omega(\phi, s)^2 \left[|\partial_\mu \tilde{L}_i|^2 + |\partial_\mu H_u|^2 \right]. \quad (3.25)$$

Therefore, using canonically-normalized fields, $\hat{\tilde{L}}_i \equiv \Omega(\phi, s) \tilde{L}_i$ and $\hat{H}_u \equiv \Omega(\phi, s) H_u$, the mass terms are rewritten as

$$\begin{aligned} \mathcal{L}_{\text{mass}}^{\tilde{L}-H_u} = & -\frac{\Omega(\phi, s)^2 \phi^2}{2} |y_{\nu 3i} \hat{\tilde{L}}_i|^2 - \frac{\Omega(\phi, s)^2 \phi^2}{2} |y_{\nu 3i}|^2 |\hat{H}_u|^2 \\ & + \left[\frac{M_{33}^* \Omega(\phi, s)^4 \phi^3}{4\sqrt{2}} y_{\nu 3i} \hat{\tilde{L}}_i \hat{H}_u + \text{h.c.} \right] \\ = & -\frac{\Omega(\phi, s)^2 \phi^2}{2} (y_\nu y_\nu^\dagger)_{33} |\hat{\tilde{L}}'_3|^2 - \frac{\Omega(\phi, s)^2 \phi^2}{2} (y_\nu y_\nu^\dagger)_{33} |\hat{H}_u|^2 \\ & + \left[\frac{M_{33}^* \Omega(\phi, s)^4 \phi^3}{4\sqrt{2}} (y_\nu y_\nu^\dagger)_{33}^{1/2} \hat{\tilde{L}}'_3 \hat{H}_u + \text{h.c.} \right], \end{aligned} \quad (3.26)$$

where we have defined $(y_\nu y_\nu^\dagger)_{33}^{1/2} \hat{\tilde{L}}'_3 \equiv y_{\nu 3i} \hat{\tilde{L}}_i$ in the second line following ref. [24]. Then, the stability condition is given by

$$|M_{33}| < 2\sqrt{2} (y_\nu y_\nu^\dagger)_{33}^{1/2} \frac{\Omega(\phi, s)^{-2}}{\phi}. \quad (3.27)$$

Using eqs. (3.12) and (3.14), it turns out that

$$|M_{33}| < 4\sqrt{2}(y_\nu y_\nu^\dagger)_{33}^{1/2} \sqrt{\beta} \\ \simeq 1.4 \times 10^{14} \text{ GeV} \left(\frac{(y_\nu y_\nu^\dagger)_{33}}{10^{-7}} \frac{\beta}{10^{-3}} \right)^{1/2}. \quad (3.28)$$

This upper bound is weaker than (3.23) in most of the parameter space, which will be seen later.

4 Neutrino mass

In this section, we derive the mass matrices for the heavy and light neutrinos. Around the global minimum, $\Omega \simeq 1$ since $\xi \ll 1$. Consequently, all the fields are canonical. Thus, the mass terms are derived similarly in global SUSY model.

4.1 Mass matrix

The superpotential (2.2) gives the Majorana masses for the light neutrinos. To see how the masses are generated, we write down the mass terms for fermionic part of N_i^c , ν_{Li} and $\tilde{S}_-$,

$$\mathcal{L}_\nu^{\text{mass}} = -\frac{1}{2}(\bar{\psi} \mathcal{M} P_L \psi + \text{h.c.}), \quad (4.1)$$

where

$$\psi = (N_1^c, N_2^c, N_3^c, \tilde{S}_-, \nu_{L1}, \nu_{L2}, \nu_{L3})^T, \quad (4.2)$$

$$\mathcal{M} = \begin{pmatrix} \tilde{M} & \tilde{m}_\nu \\ \tilde{m}_\nu^T & \mathbf{0} \end{pmatrix}. \quad (4.3)$$

$\tilde{M}$ and $\tilde{m}_\nu$ are 4×4 and 4×3 matrices, respectively, and given by

$$\tilde{M} = \begin{pmatrix} & M & \lambda_1 \langle S_+ \rangle \\ & & \lambda_2 \langle S_+ \rangle \\ & & m_\phi \\ \lambda_1 \langle S_+ \rangle & \lambda_2 \langle S_+ \rangle & m_\phi & 0 \end{pmatrix}, \quad (4.4)$$

$$\tilde{m}_\nu = \begin{pmatrix} m_\nu \\ 0 & 0 & 0 \end{pmatrix}. \quad (4.5)$$

Here $m_{\nu ij} = y_{\nu ij} \langle H_u^0 \rangle$ with $\langle H_u^0 \rangle$ being the VEV of the up-type neutral Higgs. Then mass matrix M_ν for the light neutrinos are obtained by the seesaw mechanism [17–22],

$$M_\nu = -\tilde{m}_\nu^T \tilde{M}^{-1} \tilde{m}_\nu. \quad (4.6)$$

An important consequence of the mass matrix is that the one of three light neutrinos is massless. This is because the rank of M_ν is two. Using this mass matrix, it is possible to constrain the parameters by the observed neutrino masses.

In the later discussion we assume

$$\lambda_1, \lambda_2 \ll \lambda. \quad (4.7)$$

Here, recall that there is a freedom to choose a basis for N_1^c and N_2^c . Then, M_{12} can be rotated away. As a result, M_ν is given in the following simple expression,

$$M_{\nu ij} = -\langle H_u^0 \rangle^2 \sum_{k=1}^2 \frac{y_{\nu ki} y_{\nu kj}}{M_{kk}} + \mathcal{O}\left(\frac{\lambda_{1,2}}{\lambda} \frac{m_{\nu il} m_{\nu jm}}{M_{kk}}\right). \quad (4.8)$$

Here in the second term of right-hand side, $k = 1$ or 2 and i, j, l, m can be 1–3, and we have ignored M_{i3} based on the discussion of the previous section. The leading order term, on the other hand, is independent of $\lambda_{1,2}$, M_{i3} , m_ϕ and $y_{\nu 3i}$ ($i = 1-3$). Therefore, they are not constrained by the neutrino oscillation data. This fact is important in the estimation of the reheating temperature, which we will see later. To ensure the non-zero $\lambda_{1,2}$ does not affect our later analysis at percent level, we implicitly assume $\lambda_{1,2} < 0.01\lambda$.

Before further discussing the light neutrino mass matrix, let us note that the mass matrix $\tilde{M}$ corresponds to the mass matrix in the superpotential around the global minimum,

$$W_{\text{neu}} = \frac{1}{2}(N_1^c, N_2^c, N_3^c, S_-)\tilde{M} \begin{pmatrix} N_1^c \\ N_2^c \\ N_3^c \\ S_- \end{pmatrix} + \dots \quad (4.9)$$

Recall that we have the requirement (3.23) for successful inflation. Therefore, $\tilde{M}$ should be almost block-diagonal as

$$\tilde{M} \simeq \begin{pmatrix} M_1 & & & 0 \\ & M_2 & & 0 \\ & & M_3 & m_\phi \\ 0 & 0 & m_\phi & 0 \end{pmatrix}, \quad (4.10)$$

where $M_i > 0$. Here we have left M_3 for later discussion. In the following analysis we use eq. (4.10) for $\tilde{M}$. $\tilde{M}$ is further diagonalized by a unitary matrix $U_{\tilde{M}}$ as,

$$D_{\tilde{M}} \simeq \text{diag}(M_1, M_2, m_\phi, m_\phi) \simeq U_{\tilde{M}}^T \tilde{M} U_{\tilde{M}}, \quad (4.11)$$

where

$$U_{\tilde{M}} = \begin{pmatrix} 1 & & & \\ & 1 & & \\ & & \frac{1}{\sqrt{2}} & \begin{pmatrix} 1 & i \\ 1 & -i \end{pmatrix} \end{pmatrix}. \quad (4.12)$$

Here we have omitted $\mathcal{O}(M_3/m_\phi)$ corrections since they are irrelevant in the later analysis. Non-zero $\lambda_{1,2}$ corrections enter as $\mathcal{O}(\lambda_{1,2}/\lambda)$ and $\mathcal{O}((m_\phi/M_{1,2})\lambda_{1,2}/\lambda)$ for $m_\phi \gtrsim M_{1,2}$ and $m_\phi \lesssim M_{1,2}$, respectively.

	$\Delta m_{21}^2 [10^{-5} \text{ eV}^2]$	$\Delta m_{3l}^2 [10^{-3} \text{ eV}^2]$
NH	$7.39^{+0.21}_{-0.20}$	$+2.525^{+0.033}_{-0.031}$
IH	$7.39^{+0.21}_{-0.20}$	$-2.512^{+0.034}_{-0.031}$

Table 1. Neutrino mass data taken from refs. [25, 26], adopting data with the atmospheric neutrino by Super-Kamiokande. $\Delta m_{ij}^2 \equiv m_i^2 - m_j^2$, and $\Delta m_{3l}^2 = \Delta m_{31}^2 > 0$ for the normal hierarchy (NH) case and $\Delta m_{3l}^2 = \Delta m_{32}^2 < 0$ for the inverted hierarchy (IH) case.

The corrections due to non-zero $\lambda_{1,2}$ are similar in the scalar sector since the mass matrix in $(\tilde{N}_1, \tilde{N}_2, \tilde{N}_3, S_-)$ basis is given by,

$$M_{\text{scalar}}^2 = \tilde{M}^\dagger \tilde{M}. \quad (4.13)$$

Therefore the assumption (4.7) assures that the dynamics of the inflaton and the waterfall field discussed in the previous section, as well as subsequent reheating and leptogenesis, is not affected. If it is not satisfied, then the inflationary trajectory would be altered so that we need further analysis to find the parameter space that is consistent with the CMB observations. Accordingly, the subsequent reheating and leptogenesis scenario change. We leave the detailed analysis as an interesting future work.

4.2 Parametrization of neutrino Yukawa couplings

Now let us discuss M_ν . It can be diagonalized by a unitary matrix U_ν as

$$D_{M_\nu} = \text{diag}(m_1, m_2, m_3) = U_\nu^T M_\nu U_\nu. \quad (4.14)$$

As noted above, one of the three light neutrino masses is zero. We follow the standard convention that $m_3 > m_2 > m_1 (= 0)$ for the normal hierarchy (NH) case and $m_2 > m_1 > m_3 (= 0)$ for the inverted hierarchy (IH) case and use the values given in refs. [25, 26], which are listed in table 1.

Before discussing the parametrization of the neutrino Yukawa couplings, it is instructive to count the number of parameters. The situation is the same as one discussed in refs. [24, 27] since the mass matrix of the light neutrinos (4.8) is similar. Since one neutrino is massless, there are 7 parameters in low energy, *i.e.*, 2 neutrino masses + 3 real mixing angles + 2 phases. On the other hand, M_ν includes $y_{\nu ki}$ and M_k where $k = 1, 2$ and $i = 1, 2, 3$, which means 12 real parameters (neutrino Yukawa couplings) + 2 real parameters (right-handed neutrino masses). However, 3 phases can be absorbed by lepton doublets and 2 real parameters are unphysical since M_ν is unchanged by the rescalings $y_{\nu ki} \rightarrow \gamma_k y_{\nu ki}$ and $M_k \rightarrow \gamma_k^2 M_k$ with γ_k being real constants. Therefore, we have 9 independent parameters in M_ν to determine 7 parameters in the light neutrino sector. As it will be seen below, however, the parametrization of the Yukawa couplings is different, especially for $y_{\nu 3i}$ that are important parameters for the estimation of the reheating temperature.

Let us discuss the NH case first. We define 4×3 matrix R in the similar manner in refs. [28, 29],

$$iR = D_{\tilde{M}}^{-1/2} U_{\tilde{M}}^T \tilde{m}_\nu U_\nu D_{M_\nu}^{-1/2}, \quad (4.15)$$

which satisfies

$$R^T R = \text{diag}(0, 1, 1). \quad (4.16)$$

Here $D_{\tilde{M}}^{-1/2}$ and $D_{M_\nu}^{-1/2}$ are matrices that satisfy $(D_{\tilde{M}}^{-1/2})^2 = D_{\tilde{M}}^{-1}$ and $(D_{M_\nu}^{-1/2})^2 = \text{diag}(0, m_2^{-1}, m_3^{-1})$, respectively. It is found that R is more restrictive than eq. (4.16). Namely,

$$r^T r = r r^T = \mathbf{1}, \quad (4.17)$$

$$R_{42}/R_{32} = R_{43}/R_{33} = i, \quad (4.18)$$

$$R_{l1} = 0 \quad (l = 1-4), \quad (4.19)$$

where

$$r \equiv \begin{pmatrix} R_{12} & R_{13} \\ R_{22} & R_{23} \end{pmatrix}. \quad (4.20)$$

Using the relations, the neutrino Yukawa couplings can be expressed in terms of R . For later discussion, it is useful to give following quantities:

$$(y_\nu y_\nu^\dagger)_{ii} = M_i \sum_{j=2}^3 |R_{ij}|^2 m_j / \langle H_u^0 \rangle^2 \quad (i = 1, 2), \quad (4.21)$$

$$(y_\nu y_\nu^\dagger)_{33} = 2m_\phi \sum_{j=2}^3 |R_{3j}|^2 m_j / \langle H_u^0 \rangle^2, \quad (4.22)$$

$$\text{Im} \left[(y_\nu y_\nu^\dagger)_{21}^2 \right] \frac{M_1}{M_2} = -\frac{M_1^2}{\langle H_u^0 \rangle^4} \text{Im} \left[\sum_{j=2}^3 R_{1j}^2 m_j^2 \right], \quad (4.23)$$

$$\sum_{i=1}^2 \text{Im} \left[(y_\nu y_\nu^\dagger)_{i3}^2 \right] \frac{M_3}{M_i} = -\frac{2M_3 m_\phi}{\langle H_u^0 \rangle^4} \text{Im} \left[\sum_{j=2}^3 R_{3j}^2 m_j^2 \right]. \quad (4.24)$$

It is seen that $(y_\nu y_\nu^\dagger)_{11}$ and $(y_\nu y_\nu^\dagger)_{22}$ are constrained by the neutrino oscillation data, meanwhile $(y_\nu y_\nu^\dagger)_{33}$ is basically a free parameter since R_{3j} is not constrained. This is consistent with the fact that M_ν is independent of $y_{\nu 3i}$.

The discussion is quite similar in the IH case. The definition of R is the same form as in eq. (4.15), but satisfies $R^T R = \text{diag}(1, 1, 0)$ with $(D_{M_\nu}^{-1/2})^2 = \text{diag}(m_1^{-1}, m_2^{-1}, 0)$. Then, defining r as

$$r \equiv \begin{pmatrix} R_{11} & R_{12} \\ R_{21} & R_{22} \end{pmatrix}, \quad (4.25)$$

we get

$$r^T r = r r^T = \mathbf{1}, \quad (4.26)$$

$$R_{41}/R_{31} = R_{42}/R_{32} = i, \quad (4.27)$$

$$R_{l3} = 0 \quad (l = 1-4), \quad (4.28)$$

and the neutrino Yukawa couplings are given by,

$$(y_\nu y_\nu^\dagger)_{ii} = M_i \sum_{j=1}^2 |R_{ij}|^2 m_j / \langle H_u^0 \rangle^2 \quad (i = 1, 2), \quad (4.29)$$

$$(y_\nu y_\nu^\dagger)_{33} = 2m_\phi \sum_{j=1}^2 |R_{3j}|^2 m_j / \langle H_u^0 \rangle^2, \quad (4.30)$$

$$\text{Im} \left[(y_\nu y_\nu^\dagger)_{21}^2 \right] \frac{M_1}{M_2} = -\frac{M_1^2}{\langle H_u^0 \rangle^4} \text{Im} \left[\sum_{j=1}^2 R_{1j}^2 m_j^2 \right], \quad (4.31)$$

$$\sum_{i=1}^2 \text{Im} \left[(y_\nu y_\nu^\dagger)_{i3}^2 \right] \frac{M_3}{M_i} = -\frac{2M_3 m_\phi}{\langle H_u^0 \rangle^4} \text{Im} \left[\sum_{j=1}^2 R_{3j}^2 m_j^2 \right]. \quad (4.32)$$

5 Post inflationary regime

After the end of inflation, the inflaton oscillates around the global minimum and decays eventually. Due to the decay the universe is reheated and thermal plasma is created. In this section, we estimate the reheating temperature and discuss how the lepton number asymmetry is generated. As in the previous section, we take $\Omega \simeq 1$. After the universe is reheated, gravitinos are produced in various ways. We discuss the gravitino problem at the end of this section.

5.1 Reheating

The reheating temperature T_R due to the inflaton decay is estimated by,

$$T_R \simeq (90/\pi^2 g_*(T_R))^{1/4} \sqrt{\Gamma_\phi M_{pl}}, \quad (5.1)$$

where $g_*(T)$ is the effective degree of freedom of radiation fields at temperature T and Γ_ϕ is the decay rate of the inflaton. This expression is valid when the neutrino Yukawa couplings that are responsible for the decay is sufficiently small to satisfy $T_R \lesssim m_\phi$ [30, 31], which is the situation we focus on.⁶ The inflaton decays as $\phi \rightarrow L\tilde{H}_u, \bar{L}\tilde{H}_u, \tilde{L}H_u, \tilde{L}^*H_u^*$.⁷ Here flavor indices and SU(2) doublet components are summed implicitly. Then the decay rates for the modes are given by

$$\Gamma_{\phi \rightarrow L\tilde{H}_u} = \Gamma_{\phi \rightarrow \bar{L}\tilde{H}_u} = \frac{(y_\nu y_\nu^\dagger)_{33}}{16\pi} m_\phi, \quad (5.2)$$

$$\Gamma_{\phi \rightarrow \tilde{L}H_u} = \Gamma_{\phi \rightarrow \tilde{L}^*H_u^*} = \frac{(y_\nu y_\nu^\dagger)_{33}}{16\pi} \frac{M_3^2}{m_\phi}. \quad (5.3)$$

⁶Of course, it is possible to consider a higher reheating temperature than the inflaton mass. Such a case is discussed in ref. [24]. We will comment on the impact of such high reheating temperature on leptogenesis in the next subsection.

⁷In general, the inflaton decays to gravitino pair or gravitino and right-handed neutrino. We will discuss those processes in section 5.3.

Since $\Gamma_{\phi \rightarrow \tilde{L} H_u} (= \Gamma_{\phi \rightarrow \tilde{L}^* H_u^*})$ is suppressed by $(M_3/m_\phi)^2$, the total decay rate is given by

$$\Gamma_\phi \simeq \Gamma_{\phi \rightarrow L \tilde{H}_u} + \Gamma_{\phi \rightarrow \bar{L} \tilde{\bar{H}}_u} = \frac{(y_\nu y_\nu^\dagger)_{33}}{8\pi} m_\phi. \quad (5.4)$$

Then the reheating temperature is estimated as

$$T_R \simeq 1.4 \times 10^{10} \text{ GeV} \left(\frac{m_\phi}{10^{13} \text{ GeV}} \right)^{1/2} \left(\frac{(y_\nu y_\nu^\dagger)_{33}}{10^{-9}} \right)^{1/2} \left(\frac{g_*(T_R)}{228.75} \right)^{-1/4}. \quad (5.5)$$

Recall that $(y_\nu y_\nu^\dagger)_{33}$ is not constrained by the neutrino observations. As a consequence, it is possible to consider a wide range of values for the reheating temperature, which is suitable for leptogenesis.

To end this subsection, we derive the number of e -folds before the end of inflation. In this model, the inflaton oscillates after the end of inflation and eventually decays to reheat the universe. Therefore, it is given by

$$N_e \simeq 55 + \log \left(\frac{L}{\text{Gpc}} \right) + \frac{1}{3} \log \left(\frac{T_R}{10^{10} \text{ GeV}} \right) + \log \left(\frac{H_e}{10^{13} \text{ GeV}} \right) - \frac{2}{3} \log \left(\frac{H_{\text{end}}}{10^{13} \text{ GeV}} \right), \quad (5.6)$$

where L is the present cosmological scale, and H_e and H_{end} are the Hubble scale corresponding to N_e and at the end of inflation, respectively. It is now clear that the parameters given in eq. (2.11) is consistent with $N_e = 55\text{--}60$.

5.2 Leptogenesis

Now we discuss the lepton number asymmetry. The lepton number is generated via leptogenesis [23] (see, for example, refs. [32, 33] for review). In the following numerical study, we discuss following representative cases:

$$\text{(I). } M_1, M_2 < m_\phi, \quad (5.7)$$

$$\text{(II). } M_1, M_2 > m_\phi. \quad (5.8)$$

In both cases, M_3 should satisfy eq. (3.23).

Let us consider case (I) first. For simplicity, we consider the Majorana masses are further hierarchical, i.e., $M_1 \ll M_2$.⁸ Similar situation has been studied intensively in the literature [34, 35]. Even though the lepton number is generated by the inflaton decay, it is possibly washed out when the reheating temperature is comparable or higher than M_1 .⁹

⁸ M_3 is irrelevant for leptogenesis if eq. (3.23) is satisfied.

⁹It would be possible that $\tilde{N}_i$ ($i = 1, 2$) have initial amplitude of the Hubble parameter during inflation $H_{\text{inf}} \sim g\xi/\sqrt{6}M_{pl}$. Then $\tilde{N}_i$ start to oscillate when $H \sim M_i$, and eventually decays. However, the effect of coherent oscillation of $\tilde{N}_i$ is negligible since the energy density ratio of $\tilde{N}_i$ to radiation at the decay is estimated as less than $\xi^2/18M_{pl}^4 \sim 10^{-9}$.

To see this more explicitly, it is convenient to introduce the effective neutrino mass [36] and equilibrium neutrino mass [37],

$$\tilde{m}_1 = \frac{(m_\nu m_\nu^\dagger)_{11}}{M_1}, \quad (5.9)$$

$$m_* = \frac{4\pi^2 \sqrt{g_*(M_1)} \langle H_u^0 \rangle^2}{3\sqrt{10} M_{pl}} \simeq 3.9 \times 10^{-4} \text{ eV} \left(\frac{\langle H_u^0 \rangle}{v/2} \right)^2, \quad (5.10)$$

where $v \simeq 246.7$ GeV. If $\tilde{m}_1/m_*$ is larger than unity, then it is the strong washout regime and the lepton number generated at the reheating is washed out. $\tilde{m}_1$ is estimated by using eqs. (4.21) and (4.29),

$$\tilde{m}_1 = \begin{cases} \sum_{j=2}^3 |R_{1j}|^2 m_j & \text{(NH)} \\ \sum_{j=1}^2 |R_{1j}|^2 m_j & \text{(IH)} \end{cases}. \quad (5.11)$$

Using the neutrino mass data and eqs. (4.17) and (4.26), it is straightforward to find that $\tilde{m}_1$ has a lower bound,

$$\tilde{m}_1 \geq \begin{cases} m_2 \simeq 8.6 \times 10^{-3} \text{ eV} & \text{(NH)} \\ m_1 \simeq 4.9 \times 10^{-2} \text{ eV} & \text{(IH)} \end{cases}. \quad (5.12)$$

Therefore, it is the strong washout regime in either case.

Although the primordial lepton number is washed out, the lepton number is regenerated by the decay of the lightest right-handed (s)neutrino, i.e., N_1 and $\tilde{N}_1$ in the present case. Then the lepton number, strictly speaking lepton number minus baryon number, is converted to baryon number via the sphaleron effect. This scenario works if $T_R \gtrsim M_1$ [34, 35], which is always possible as confirmed in the previous subsection. Then the resultant baryon number becomes independent of T_R . In our study, we adopt the analytic expressions in ref. [35] for the calculation of the baryon number. Note that although the results there are given in non-supersymmetric model, the results in supersymmetric model do not change much both quantitatively and qualitatively [33, 38, 39]. In our study we adopt the discussion given in ref. [33]. Then the baryon number is determined by

$$\eta_B \equiv \frac{n_B}{n_\gamma} = \frac{3\sqrt{2}}{4} \frac{a_{\text{sph}}}{f} \epsilon_1 \kappa_f \simeq 2.7 \times 10^{-10} \left(\frac{\epsilon_1}{10^{-6}} \right) \left(\frac{\kappa_f}{2 \times 10^{-2}} \right), \quad (5.13)$$

where n_B and n_γ are number densities of baryon and photon at present, respectively, $a_{\text{sph}} = 28/79$, $f = 2387/86$, and a factor of $\sqrt{2}$ counts the supersymmetric effect. The efficiency factor κ_f is given by [35]

$$\kappa_f = (2 \pm 1) \times 10^{-2} \left(\frac{0.01 \text{ eV}}{\tilde{m}_1} \right)^{1.1 \pm 0.1}. \quad (5.14)$$

Finally, referring ref. [40], the asymmetric parameter ϵ_1 in our model is given by

$$\epsilon_1 = -\frac{3}{16\pi} \frac{\text{Im} \left[(y_\nu y_\nu^\dagger)_{21}^2 \right]}{(y_\nu y_\nu^\dagger)_{11}} \frac{M_1}{M_2} = \frac{3}{16\pi} \frac{M_1}{\langle H_u^0 \rangle^2} m_{\text{eff}}, \quad (5.15)$$

Figure 1. Allowed region for M_1 as function of effective neutrino mass $\tilde{m}_1$ defined in eq. (5.9) for the normal hierarchy (NH) and inverted hierarchy (IH) cases. Lower bound on M_1 is obtained from $\eta_B^{\max} \geq \eta_B^{\text{obs}}$ while upper bound is (5.7). Lower bound on $\tilde{m}_1$ is given by eq. (5.12).

where we have used eqs. (4.21), (4.23), (4.29), and (4.31) to obtain

$$m_{\text{eff}} = \begin{cases} \frac{\text{Im} \sum_{j=2}^3 R_{1j}^2 m_j^2}{\sum_{j=2}^3 |R_{1j}|^2 m_j} & (\text{NH}) \\ \frac{\text{Im} \sum_{j=1}^2 R_{1j}^2 m_j^2}{\sum_{j=1}^2 |R_{1j}|^2 m_j} & (\text{IH}) \end{cases}. \quad (5.16)$$

It turns out that the maximum value of m_{eff} , denoted as $m_{\text{eff}}^{\max}$, is

$$m_{\text{eff}}^{\max} = \begin{cases} m_3 - m_2 \simeq 4.2 \times 10^{-2} \text{ eV} & (\text{NH}) \\ m_2 - m_1 \simeq 7.4 \times 10^{-4} \text{ eV} & (\text{IH}) \end{cases}. \quad (5.17)$$

Therefore, parametrizing m_{eff} as $m_{\text{eff}} = m_{\text{eff}}^{\max} \sin \delta$, the asymmetric parameter is given by

$$\epsilon_1 \simeq \begin{cases} 8.2 \times 10^{-7} & (\text{NH}) \\ 1.5 \times 10^{-8} & (\text{IH}) \end{cases} \times \left(\frac{M_1}{10^{10} \text{ GeV}} \right) \left(\frac{\langle H_u^0 \rangle}{v/2} \right)^{-2} \left(\frac{\sin \delta}{0.5} \right). \quad (5.18)$$

Using eqs. (5.13), (5.14) and (5.18), the lower limit on M_1 to explain the present baryon number is determined from $\eta_B^{\max} \geq \eta_B^{\text{obs}}$ where $\sin \delta = 1$ and η_B^{obs} is given by [41]

$$\eta_B^{\text{obs}} = (6.12 \pm 0.03) \times 10^{-10}. \quad (5.19)$$

In figure 1, allowed regions are depicted for the NH and IH cases. Here we consider so-called high-scale SUSY and take $\langle H_u^0 \rangle = v/2$ to get 125 GeV Higgs mass [42, 43]. In the plot upper bound on M_1 is given by (5.7), i.e., $M_1 < m_\phi = 10^{13} \text{ GeV}$, and the lower bound on

$\tilde{m}_1$ is from eq. (5.12).¹⁰ The theoretical uncertainties in eq. (5.14) are taken into account. It is found that the present baryon number can be explained in a wide range of parameter space for the NH case. For the IH case, on the other hand, it seems that the parameter space for a successful leptogenesis is relatively limited. The lowest value required for M_1 turns out to be

$$M_1 \gtrsim \begin{cases} 5.8 \times 10^9 \text{ GeV} & (\text{NH}) \\ 2.1 \times 10^{12} \text{ GeV} & (\text{IH}) \end{cases} . \quad (5.20)$$

The lower limit is near the upper bound in the IH case. Here recall that the upper bound on M_1 is just a theoretical one. When $M_1 \sim m_\phi$, T_R should be comparable to m_ϕ , which is possible as discussed in refs. [24, 30, 31]. In such a case, sneutrino inflation and leptogenesis can be another source for lepton asymmetry, which will be discussed below in detail. Therefore, the upper bound merely indicates the parameter space for simple thermal leptogenesis to work.

Let us move on to case (II). Since they are much heavier than the inflaton, $N_{1,2}$ and $\tilde{N}_{1,2}$ are never thermalized after the reheating. For N_3 and $\tilde{N}_3$, on the other hand, it depends on the effective neutrino mass that is defined by

$$\tilde{m}_3 = \frac{(m_\nu m_\nu^\dagger)_{33}}{m_\phi} \simeq 1.5 \times 10^{-9} \text{ eV} \left(\frac{(y_\nu y_\nu^\dagger)_{33}}{10^{-9}} \right) \left(\frac{10^{13} \text{ GeV}}{m_\phi} \right) \left(\frac{\langle H_u^0 \rangle}{v/2} \right)^2 . \quad (5.21)$$

Then, from eq. (5.5), one obtains

$$\frac{T_R}{m_\phi} \simeq 0.71 \times \left(\frac{\tilde{m}_3}{m_*} \right)^{1/2} . \quad (5.22)$$

Here m_* is defined similarly to eq. (5.10) but replacing $g_*(M_1)$ by $g_*(m_\phi)$, and we have taken $g_*(T_R) \simeq g_*(m_\phi)$. As explained in section 5.1, the expression eq. (5.5) is valid for $T_R/m_\phi \lesssim 1$ that is satisfied for $\tilde{m}_3 < m_*$. Such a case corresponds to the weak washout regime. In that regime, the N_3 and $\tilde{N}_3$ are not thermalized, and the lepton number produced by the inflaton decay can be the source of the present baryon number. This situation is similar to sneutrino inflation and leptogenesis [24, 27, 44–52]. (See also refs. [53–55] for leptogenesis via Afflec-Dine mechanism [56].)

Let us suppose $\tilde{m}_3 \ll m_*$, i.e., $T_R/m_\phi \ll 1$. Then baryon number is given by

$$\eta_B = \frac{3}{4} \frac{T_R}{m_\phi} a_{\text{sph}}^{\text{MSSM}} d \epsilon_\phi , \quad (5.23)$$

where $a_{\text{sph}}^{\text{MSSM}} = 8/23$ and $d = (s/n_\gamma)_0 = 43\pi^4/495\zeta(3)$ is the present value of entropy density to and photon density ratio. ϵ_ϕ is obtained by an explicit calculation as

$$\epsilon_\phi = -\frac{3}{4\pi} \sum_{i=1}^2 \frac{\text{Im} \left[(y_\nu y_\nu^\dagger)_{i3}^2 \right]}{(y_\nu y_\nu^\dagger)_{33}} \frac{M_3}{M_i} = \frac{3}{4\pi} \frac{M_3}{\langle H_u^0 \rangle^2} m'_{\text{eff}} , \quad (5.24)$$

¹⁰Strictly speaking, the equality should be excluded since baryon number is zero.

where

$$m'_{\text{eff}} = \begin{cases} \frac{\text{Im} \sum_{j=2}^3 R_{3j}^2 m_j^2}{\sum_{j=2}^3 |R_{3j}|^2 m_j} & (\text{NH}) \\ \frac{\text{Im} \sum_{j=1}^2 R_{3j}^2 m_j^2}{\sum_{j=1}^2 |R_{3j}|^2 m_j} & (\text{IH}) \end{cases} . \quad (5.25)$$

In the second step, we have used eqs. (4.22), (4.24), (4.30), and (4.32). It should be noted that ϵ_ϕ is independent of the inflaton mass, but it depends on M_3 . Even though R_{3j} are not constrained, it has been found that m'_{eff} is bounded from above. The maximum value turns out to be

$$m'_{\text{eff}}^{\text{max}} = \begin{cases} m_3 \simeq 5.0 \times 10^{-2} \text{ eV} & (\text{NH}) \\ m_2 \simeq 5.0 \times 10^{-2} \text{ eV} & (\text{IH}) \end{cases} . \quad (5.26)$$

Therefore, parametrizing the effective neutrino mass as $m'_{\text{eff}} = m'_{\text{eff}}^{\text{max}} \sin \delta'$ the asymmetric parameter is given by

$$\epsilon_\phi \simeq 3.9 \times 10^{-9} \times \left(\frac{M_3}{10^7 \text{ GeV}} \right) \left(\frac{\langle H_u^0 \rangle}{v/2} \right)^{-2} \left(\frac{\sin \delta'}{0.5} \right) , \quad (5.27)$$

for both the NH and IH cases. Using the equations, we get

$$\eta_B^{\text{max}} \simeq 1.6 \times 10^{-10} \left(\frac{M_3}{10^7 \text{ GeV}} \right) \left(\frac{\tilde{m}_3}{10^{-7} \text{ eV}} \right)^{1/2} \left(\frac{\langle H_u^0 \rangle}{v/2} \right)^{-3} . \quad (5.28)$$

Since η_B is independent of m_ϕ , the requirement $\eta_B^{\text{max}} \geq \eta_B^{\text{obs}}$ gives a lower bound on M_3 , which is plotted in figure 2. Here $\langle H_u^0 \rangle = v/2$ is taken as in figure 1. Upper bound $M_3 < 2 \times 10^{11} \text{ GeV}$ is from eq. (3.23). Another upper bound on M_3 from the stability of the inflationary trajectory, eq. (3.28), is also shown. We quit plotting for region $\tilde{m}_3 > 10^{-5} \text{ eV}$ because $\tilde{m}_3 \ll m_*$ is no longer valid. In the plot, contours of T_R are depicted. It is found that the leptogenesis is successful in a wide parameter space for both the NH and IH cases. Lower bound on M_3 behaves similarly to region C in figure 1 of ref. [48] by reading M_1 and $\tilde{m}_1$ as M_3 and $\tilde{m}_3$, respectively, i.e., the lower bound is proportional to $1/\sqrt{\tilde{m}_3}$. Quantitatively, the lower bound in our model is relaxed by roughly a factor of 4 compared to the result in the reference. This can be understood as follows; first, the decay rate of the inflaton in our model is different from one in the literature (see eq. (5.4)), leading to a $1/\sqrt{2}$ suppression in T_R/m_ϕ ; as another consequence, the expression (5.24) (with $m'_{\text{eff}} = m'_{\text{eff}}^{\text{max}}$) is enhanced by a factor of two compared to $|\epsilon_1^{\text{max}}|$ in the reference; finally $\tan \beta = \infty$ is taken in the work meanwhile $\tan \beta = 1$ in our model. Thus, in total, a factor of 4 enhancement is obtained in η_B .

In the case where $\tilde{m}_3$ gets much larger than m_* , the situation reduces to case (I). Namely, the reheating temperature is so high that both $\tilde{N}_3$ and N_3 are thermalized and thermal leptogenesis takes place. Resultant allowed region is the same as the NH case of figure 1, by replacing M_1 and $\tilde{m}_1$ by M_3 and $\tilde{m}_3$, respectively, but there is no lower bound on $\tilde{m}_3$ meanwhile there is the upper bound on M_3 . In the intermediate case, $\tilde{m}_3 \sim m_*$, on the other hand, the Boltzmann equations should be solved numerically to get the lepton

Figure 2. Allowed region for M_3 as function of $\tilde{m}_3$ defined in eq. (5.21). Results are the same for the NH and IH cases. Lower bound on M_3 comes from $\eta_B^{\max} \geq \eta_B^{\text{obs}}$ meanwhile upper one is given by eqs. (3.23) and (3.28). The former is not to affect the inflationary trajectory and the latter (dashed orange) is from the stability of the scalar potential. Vertical lines (dot-dashed purple) are contours of T_R .

number, which is already done in ref. [48]. The result corresponds to region B in figure 1 of the reference. Strictly speaking, the effective dissipation rate should be used instead of the decay rate of the inflaton [31] in the Boltzmann equations. As shown in the reference, the reheating process is so efficient when the dissipation rate is taken into account that the reheating temperature can exceed the mass of the inflaton mass and consequently N_3 and $\tilde{N}_3$ are easily thermalized. Once they are thermalized, the thermal leptogenesis takes place, where the resultant baryon number becomes independent of the reheating temperature. Eventually the situation reduces to the case (I). Such qualitative behavior can be confirmed by numerical study, which is left for the future work.

Crucial difference from sneutrino leptogenesis [44–48] is that although M_3 and $\tilde{m}_3$, i.e., $(y_\nu y_\nu^\dagger)_{33}$, are important parameters to determine baryon number, they are sequestered from other physical quantities, such as the heavy right-handed (s)neutrino masses or the light neutrino mass matrix. Therefore, there is no consequence in other low energy experiments. This is a feature of case (II).

5.3 Gravitino problem

In the framework of supergravity, a fair amount of gravitino ψ_μ can be produced in various ways in the thermal history of the universe. Since the interactions of gravitino with the MSSM particles are Planck-suppressed, gravitino is long-lived and its decay can spoil the successful big-bang nucleosynthesis if it is unstable. Although this problem can be avoided

when gravitino is enough heavy to have the lifetime much shorter than 1 sec, gravitino decay produces the lightest superparticle (LSP). Then the LSP produced by the decay may overclose the universe if the R-parity is conserved.

There are three types of production mechanism of gravitino in the model we consider; (i) the inflaton decay; (ii) thermal scattering from the thermal bath [57–60]; (iii) decay of superparticles in the thermal bath [61, 62].

In general, process (i) includes gravitino pair production. The decay width of the mode, however, depends on the inflaton VEV [63–68]. In our case, therefore, this process can be ignored since the inflaton does not have a VEV. On the other hand, the inflaton can decay to gravitino and right-handed neutrino. The decay width is given by

$$\Gamma_{\phi \rightarrow \psi_\mu N_3} = \frac{\beta_f^3 m_\phi^5}{48\pi M_{pl}^2 m_{3/2}^2} \left[1 - \frac{(m_f + m_{3/2})^2}{m_\phi^2} \right], \quad (5.29)$$

where $\beta_f^2 = 1 - 2(m_f^2 + m_{3/2}^2)/m_\phi^2 + (m_f^2 - m_{3/2}^2)^2/m_\phi^4$. Here m_f and $m_{3/2}$ are masses of N_3 and ψ_μ , respectively. The mass difference between ϕ and N_3 is expected to be given by the soft SUSY breaking mass scale for scalar superpartners, $|m_\phi - m_f| \sim \tilde{m}$. Let us assume this decay happens by taking $m_\phi = m_f + \tilde{m}$ where $\tilde{m} = km_{3/2}$ ($k > 1$). In the limit $km_{3/2}/m_\phi \ll 1$, we obtain

$$\Gamma_{\phi \rightarrow \psi_\mu N_3} \simeq \frac{C m_\phi m_{3/2}^2}{3\pi M_{pl}^2}, \quad (5.30)$$

where $C = (k-1)(k^2-1)^{3/2}$. Then the branching fraction of this mode is

$$\text{Br}_{\phi \rightarrow \psi_\mu N_3} = \frac{\Gamma_{\phi \rightarrow \psi_\mu N_3}}{\Gamma_\phi} \simeq 4.5 \times 10^{-16} C \left(\frac{m_{3/2}}{10^5 \text{ GeV}} \right)^2 \left(\frac{10^{-11}}{(y_\nu y_\nu^\dagger)_{33}} \right). \quad (5.31)$$

Therefore, it is sure that the inflaton decay reheats the universe in wide range of gravitino mass region. On the other hand, the resultant gravitino abundance produced by the decay is estimated as

$$\Omega_{3/2}^{\text{inf}} h^2 \sim 1.3 \times 10^{-6} C \left(\frac{m_{3/2}}{10^5 \text{ GeV}} \right)^3 \left(\frac{10^{13} \text{ GeV}}{m_\phi} \right)^{1/2} \left(\frac{10^{-11}}{(y_\nu y_\nu^\dagger)_{33}} \right)^{1/2}, \quad (5.32)$$

where h is the scale factor of Hubble expansion rate.

Gravitino abundance via process (ii) is most effective at high temperature, thus it is proportional to T_R , meanwhile in process (iii) gravitino is dominantly produced when the temperature is around the mass of decaying particle. Adopting the expression given in ref. [69], the abundances via processes (ii) and (iii) are given by

$$\Omega_{3/2}^{\text{TH}} h^2 \sim 4 \left(\frac{T_R}{10^9 \text{ GeV}} \right) \left(\frac{m_{3/2}}{10^5 \text{ GeV}} \right), \quad (5.33)$$

$$\Omega_{3/2}^{\text{FI}} h^2 \sim 5 \times 10^{-4} k^3 \left(\frac{m_{3/2}}{10^5 \text{ GeV}} \right)^2. \quad (5.34)$$

Figure 3. Relic density of the LSP as function of gravitino mass. “inf” (dotted green), “TH” (dot-dashed red), and “FI” (dashed blue) are the contributions from process (i), (ii), and (iii), respectively. $\Omega_{\text{DM}} h^2 \simeq 0.12$ (solid brown) is indicated as a reference. Right region from vertical line (dotted violet) indicates $T_{3/2} > m_{\text{LSP}}$, and shaded region is excluded. $m_{\text{LSP}} = 1$ TeV, $m_\phi = 10^{13}$ GeV, $(y_\nu y_\nu^\dagger)_{33} = 10^{-11}$, and $k = \tilde{m}/m_{3/2} = 2$ (left), 10 (right) are taken.

Here the contribution of the longitudinal mode of gravitino is suppressed in $\Omega_{3/2}^{\text{TH}}$ by considering gluino is lighter than gravitino. In $\Omega_{3/2}^{\text{FI}}$, we have assumed that all scalar leptons and quarks in the MSSM (whose mass scale is $\tilde{m}$) are heavier than gravitino and that they are thermalized. And as in the discussion of the inflaton decay, $\tilde{m} = k m_{3/2}$ has been taken.¹¹ Then the relic abundance of the LSP due to gravitino decay is estimated as

$$\Omega_{\text{LSP}}^{\text{non-th}} h^2 = \frac{m_{\text{LSP}}}{m_{3/2}} \Omega_{3/2} h^2, \quad (5.35)$$

where m_{LSP} is the LSP mass and $\Omega_{3/2} = \Omega_{3/2}^{\text{inf}} + \Omega_{3/2}^{\text{TH}} + \Omega_{3/2}^{\text{FI}}$ is the total gravitino abundance. $\Omega_{\text{LSP}}^{\text{non-th}} h^2$ should not exceed the observed dark matter abundance $\Omega_{\text{DM}} h^2 \simeq 0.12$ [41], which gives a constraint on gravitino mass.

Figure 3 shows the resultant $\Omega_{\text{LSP}}^{\text{non-th}}$ from the each contribution. Here $m_{\text{LSP}} = 1$ TeV is taken by considering 1 TeV Higgsino or Wino dark matter with a mass of 2.7–3 TeV [70]. $m_\phi = 10^{13}$ GeV, $(y_\nu y_\nu^\dagger)_{33} = 10^{-11}$, and $k = 2$ (left), 10 (right) are taken to determine the contributions from processes (i) and (iii). It is seen that in lower gravitino mass region, the dominant contribution to $\Omega_{\text{LSP}}^{\text{non-th}}$ is from process (ii). In order for the contribution not to exceed the dark matter abundance, $T_R \lesssim 10^9$ GeV is required, which is well-known result. This gives a stringent constraint on the parameter space for successful leptogenesis shown in figures 1 and 2 if gravitino is not heavy enough. Namely, $\tilde{m}_1 \sim 10^{-2}$ eV is the allowed region in case (I) meanwhile 10^{-13} eV $\lesssim \tilde{m}_3 \lesssim 10^{-11}$ eV is allowed in case (II). On the other hand, the contributions from processes (i) and (iii) depend on the parameters, especially $k = \tilde{m}/m_{3/2}$ (and $m_{3/2}$). In order for $\Omega_{\text{LSP}}^{\text{non-th}}$ not to exceed the dark matter abundance, the upper bound on gravitino mass is obtained depending on the mass spectrum of squarks and sleptons, e.g., $m_{3/2} \lesssim 10^8$ (10^6) GeV for $\tilde{m} \sim m_{3/2}$ ($10 m_{3/2}$). Such gravitino mass is

¹¹We have checked that the contribution from $\tilde{N}_1$ is negligible even if $\tilde{N}_1$ is thermalized, i.e., in case (I).

preferred in minimal or mini split supersymmetry [71, 72], pure gravity mediation [73, 74], and spread supersymmetry [69, 75].

On the other hand, there is also an allowed region in higher gravitino mass region. This is because gravitino decays before thermal freeze-out of the LSP in that region. The allowed region can be estimated by imposing gravitino decay temperature $T_{3/2}$ larger than the LSP mass. The gravitino decay temperature is defined by

$$T_{3/2} \simeq (90/\pi^2 g_*(T_{3/2}))^{1/4} \sqrt{\Gamma_{3/2} M_{pl}}, \quad (5.36)$$

where $\Gamma_{3/2}$ is the decay rate of gravitino. Then $m_{3/2} \gtrsim 3 \times 10^8 \text{ GeV}$ is obtained from $T_{3/2} > m_{\text{LSP}}$ for $m_{\text{LSP}} = 1 \text{ TeV}$ (see e.g., ref. [76]). Such high gravitino mass can be considered high-scale SUSY [77], intermediate scale supersymmetry [78], and unified inflation model [16]. However, gravitino cannot be too heavy because $\text{Br}_{\phi \rightarrow \psi_\mu N_3}$ should be less than unity for the reheating. Taking $\text{Br}_{\phi \rightarrow \psi_\mu N_3} \lesssim 0.1$, for example, upper bound on gravitino mass is obtained as $m_{3/2} \lesssim 5 \times 10^{11} \text{ GeV}$ ($2 \times 10^{10} \text{ GeV}$) for $k = 2$ (10).

Another option is the R-parity violation. Under the R-parity violation, the LSP decays to the standard-model particles. Then, the LSP does not contribute to the matter abundance of the universe so that there is no constraint on $m_{3/2}$.

6 Conclusion

Superconformal subcritical hybrid inflation is one of attractive inflation models that are consistent with the observed cosmological parameters by the Planck satellite. In this paper we have studied the cosmology of an extended version of the model. In the model three right-handed neutrinos are introduced. The superpotential consists of one in the supersymmetric seesaw model and the interaction terms of the right-handed neutrinos with the additional matter fields, one of which plays the role of the waterfall field. In the Kähler potential, on the other hand, it is possible for the sneutrinos to have shift symmetry by introducing explicit superconformal breaking terms of $\mathcal{O}(1)$. Due to the shift symmetry, one of the sneutrinos becomes the inflaton field similarly in superconformal subcritical hybrid inflation. Although the mass terms of the sneutrinos can affect the trajectory of the inflaton, it has turned out that the effect is restrictive and viable inflation is realized. After inflation, the inflaton field decays to Higgses and sleptons or Higgsinos and leptons to reheat the universe.

Light neutrino masses are given by the seesaw mechanism. However, the mass matrix is different from the conventional one. It turns out that one of the neutrinos is massless. Assuming that suppressed couplings of the other right-handed neutrinos to the waterfall field, it has been found that the neutrino Yukawa couplings that couples the inflaton to the MSSM sector are not constrained by the neutrino oscillation data. Consequently, the reheating temperature is a free parameter, which is suitable for leptogenesis.

We have considered two representative cases; (I) the other right-handed (s)neutrinos are lighter than the inflaton; (II) the other right-handed (s)neutrinos are heavier than the inflaton. In case (I), thermal leptogenesis is possible if the reheating temperature is larger

than $\sim 10^9$ GeV. It has been found leptogenesis is successful in a wide range of parameter space in the normal hierarchy case while the parameter space for leptogenesis is relatively limited in the inverted hierarchy case. In case (II), on the other hand, sneutrino leptogenesis takes place if the reheating temperature is larger than $\sim 10^8$ GeV. It has turned out that in both the normal and inverted hierarchy cases successful leptogenesis is realized in wide range of parameter space.

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