arXiv:1601.00192v1 [hep-ph] 2 Jan 2016

UT-15-42, TU-1014

Minimal Sneutrino Chaotic Inflation

Kazunori Nakayama[a,b], Fuminobu Takahashi[c,b] and Tsutomu T. Yanagida[b,d]

[a] *Department of Physics, Faculty of Science, The University of Tokyo, Bunkyo-ku, Tokyo 133-0033, Japan*

[b] *Kavli IPMU (WPI), UTIAS, The University of Tokyo, Kashiwa, Chiba 277-8583, Japan*

[c] *Department of Physics, Tohoku University, Sendai, Miyagi 980-8578, Japan*

[d] *ARC Centre of Excellence for Particle Physics at the Terascale, School of Physics, The University of Melbourne, Victoria 3010, Australia*

Abstract

We study a sneutrino chaotic inflation model with only two right-handed neutrino superfields, where one of them plays the role of the inflaton, while the other is necessary to stabilize the inflaton potential. We assume that the shift symmetry of the inflaton is explicitly broken down to discrete one in the superpotential. As a result, the neutrino Yukawa couplings are periodic in the inflaton field, and the masses of leptons and Higgs do not blow up even if the inflaton takes super-Planckian field values. The inflaton potential is given by a sum of sinusoidal functions with different height and periodicity, the so-called multi-natural inflation. We show that the predicted scalar spectral index and tensor-to-scalar ratio lie in the region favored by the Planck data. The reheating temperature is considered to be so high that successful thermal leptogenesis is possible.

1 Introduction

Cosmic microwave background (CMB) temperature and polarization anisotropies have coherence beyond the horizon at the last scattering. This clearly shows that our Universe has experienced accelerated expansion, i.e., inflation, at a very early stage of the evolution. In particular, a single-field slow-roll inflation is consistent with the observations.

Among many inflation models so far, there is an interesting class of models called large-field or chaotic inflation [1]. One of the advantages of the chaotic inflation is that it has no initial condition problem. With the chaotic initial condition at the Planckian epoch, some patch of the Universe will necessarily start to inflate. For this, the initial inflaton field value has to be larger than the Planck scale $M_P(\simeq 2.4 \times 10^{18}\,\mathrm{GeV})$ by many orders of magnitude, and it is customary to impose a shift symmetry on the inflaton to keep the inflaton potential under control.

The simplest chaotic inflation is based on the quadratic potential where the inflaton mass is fixed to be order $10^{13}\,\mathrm{GeV}$ by the normalization of the curvature perturbations. The suggested inflaton mass is intriguingly close to the right-handed (RH) neutrino mass scale required by the seesaw mechanism [2] for the light neutrino masses, making it plausible to identify the inflaton with the supersymmetric (SUSY) partner of the RH neutrino. Another virtue of the sneutrino chaotic inflation is that the inflaton automatically reheats the minimal supersymmetric standard model (MSSM) particles via the decay into leptons and Higgs. As a result, thermal (or non-thermal) leptogenesis [3] takes places naturally.

A first attempt to build a sneutrino chaotic inflation model in supergravity relied on a rather complicated form of the Kähler potential [4, 5]. More recently, a much simpler realization was proposed [6], based on a generic construction of chaotic inflation models in supergravity [7] (see also [8]). The recent Planck observations [9], however, excluded the quadratic chaotic inflation model, which requires some modifications of the inflaton potential such as the polynomial chaotic inflation [10–12].

The purpose of the present paper is twofold. First, we study a minimal sneutrino chaotic inflation with only two RH neutrinos, $N_{1,2}$, where one of them (N_1) plays the role of the inflaton while the other (N_2) is necessary to ensure the stability of the inflaton potential. The latter is often called the stabilizer field. Interestingly, the neutrino oscillation data can be explained with only two RH neutrinos [13], and the simple realization of the sneutrino chaotic inflation also requires at least two RH neutrinos. In this minimal framework with the seesaw mechanism, the inflaton as well as the stabilizer field have sizable neutrino Yukawa couplings of $\mathcal{O}(0.1)$, and therefore, the chaotic initial condition with super-Planckian inflaton field values cannot be realized. This is because the masses of the leptons and Higgs would exceed the Planck mass for the inflaton field value greater than $\mathcal{O}(10)M_P$, and then, the effective field theory description breaks down. Our second purpose is to propose a solution to the problem and study its implications. In fact, the problem can be solved if the neutrino Yukawa couplings are some functions of the inflaton so that the masses of the leptons and Higgs do not monotonically increase with the inflaton field; for example, they may asymptote to a constant value or start to decrease for super-Planckian inflaton field values. An interesting

possibility is that the system comes back to the SM as the inflaton field exceeds a critical value. This is the case if the shift symmetry of the inflaton is not completely broken by the superpotential interactions, but there remains an unbroken discrete shift symmetry. That is to say, the superpotential interactions are invariant under the following discrete shift transformation along the imaginary direction,

$$\varphi \to \varphi + 2\pi f, \tag{1.1}$$

where $\varphi \equiv \sqrt{2}\mathrm{Im}[N_1]$, and f is the decay constant. In general, the inflaton potential is given by a sum of sinusoidal functions with different height and potential, the so-called multi-natural inflation [14–16].[#1] To avoid the blow-up of the masses of leptons and Higgs, f should be smaller than or comparable to $\mathcal{O}(10)M_P$. Thus, the prediction of the scalar spectral index (n_s) as well as the tensor-to-scalar ratio (r) are naturally deviated from the simple quadratic chaotic inflation in the minimal sneutrino inflation. As we shall see, if $f = \mathcal{O}(10)M_P$, the predicted (n_s, r) can lie in the range preferred by the Planck data.

2 Minimal sneutrino chaotic inflation

The model we consider includes the MSSM fields plus two RH neutrinos N_1 and N_2, and the Kähler and super-potentials relevant for the inflation are

$$K = \frac{1}{2}(N_1 + N_1^\dagger)^2 + |N_2|^2 - k_2\frac{|N_2|^4}{M_P^2}, \tag{2.1}$$

$$W = MN_2\sum_{n=1}\frac{g_n}{2n}f\left(e^{\frac{\sqrt{2}nN_1}{f}} - e^{\frac{-\sqrt{2}nN_1}{f}}\right) + \left[y_{1\alpha}\sum_{n=1}\frac{g'_n}{2n}f\left(e^{\frac{\sqrt{2}nN_1}{f}} - e^{\frac{-\sqrt{2}nN_1}{f}}\right) + y_{2\alpha}N_2\right]L_\alpha H_u, \tag{2.2}$$

where α runs over the lepton flavor e, μ, and τ, k_2 is a positive constant of order unity, and we take $g_1 = g'_1 = 1$. Here and in what follows, the summation over repeated indices is understood. The above Kähler potential respects a continuous shift symmetry along the imaginary direction of N_1, which ensures the flatness of the inflaton potential at $\varphi \gtrsim M_P$. The shift symmetry is explicitly broken down to a discrete one (1.1) by the above superpotential interactions. We also impose a Z_2 symmetry under which both N_1 and N_2 flip the sign.[#2] Here we take a basis of L_α such that charged lepton yukawa sector is diagonalized and omitted other interactions of the MSSM fields. The RH neutrino mass parameter M can be taken real and positive without loss of generality. The coefficients g_n and g'_n are assumed to be suppressed for larger n.

[#1] A sizable running spectral index can be generated in large field inflation with modulations [17,18]. See also Refs. [19–21].

[#2] This Z_2 can be identified with $Z_2^{(\mathrm{B-L})}$ under which all quarks and leptons flip the sign.

For $|N_1| \ll f$, the model is reduced to

$$K = \frac{1}{2}(N_1 + N_1^\dagger)^2 + |N_2|^2 - k_2 \frac{|N_2|^4}{M_P^2}, \tag{2.3}$$

$$W = M N_1 N_2 + y_{i\alpha} N_i L_\alpha H_u, \tag{2.4}$$

where $i = 1, 2$. Since typical values of $y_{i\alpha}$ is of $\mathcal{O}(0.1)$ for reproducing the observed neutrino masses (see Sec. 3.1), the masses of $L_\alpha H_u$ would exceed the Planck mass during inflation if this effective theory holds up to $|\mathrm{Im} N_1| \sim \mathcal{O}(10) M_P$. However, thanks to the discrete shift symmetry, the actual superpotential is given by (2.2), where the masses of $L_\alpha H_u$ are periodic with respect to $\varphi (= \sqrt{2}\mathrm{Im} N_1)$. Thus their masses remain smaller than the Planck mass, in which case we can safely discuss the inflaton dynamics.

2.1 Inflaton potential

In our model, the inflaton φ is identified with the imaginary component of N_1, $\varphi = \sqrt{2}\mathrm{Im}[N_1]$. The inflaton potential is then given by

$$V = M^2 \left| \sum_n \frac{g_n}{n} \sin\left(\frac{n\varphi}{f}\right) \right|^2, \tag{2.5}$$

where we have taken $\mathrm{Re}[N_1] = N_2 = H_u = L_\alpha = 0$. To be precise, $\mathrm{Re}[N_1]$ slightly shifts from zero, but its effect on the inflationary prediction is negligible [12]. We will check the validity of this assumption later. The inflaton potential is given by a sum of sinusoidal functions with different height and periodicity, and it is the so-called multi-natural inflation [14].

To proceed, let us take the first two terms:

$$\begin{aligned} V &= \frac{1}{2} M^2 f^2 \left| \sin\left(\frac{\varphi}{f}\right) + \frac{g_2}{2} \sin\left(\frac{2\varphi}{f}\right) \right|^2, \\ &= \frac{1}{2} M^2 f^2 \sin^2\left(\frac{\varphi}{f}\right) \left[1 + 2C \cos\theta \cos\left(\frac{\varphi}{f}\right) + C^2 \cos^2\left(\frac{\varphi}{f}\right) \right], \end{aligned} \tag{2.6}$$

where we have used $g_1 = 1$ and defined $g_2 \equiv C e^{i\theta}$. For $C = 0$ this is nothing but a potential for natural inflation, but a nonzero C deforms the inflaton potential and the prediction of the spectral index n_s and the tensor-to-scalar ratio r are modified. In Fig. 1 we show the shape of the scalar potential (2.6) for $C = 0.9$ with $\theta = 0, \pi/2, 9\pi/16$, and $3\pi/4$ for fixed f and M. The case of $C = 0$ is also shown for comparison. We have numerically checked that there are no local minima of the potential in the field range $0 < \varphi/f < \pi$ for $0 \leq C < 1$.

We have numerically solved the slow-roll equation of motion

$$\ddot{\varphi} + 3H\dot{\varphi} + V' = 0, \tag{2.7}$$

$$3M_P^2 H^2 \simeq V(\varphi), \tag{2.8}$$

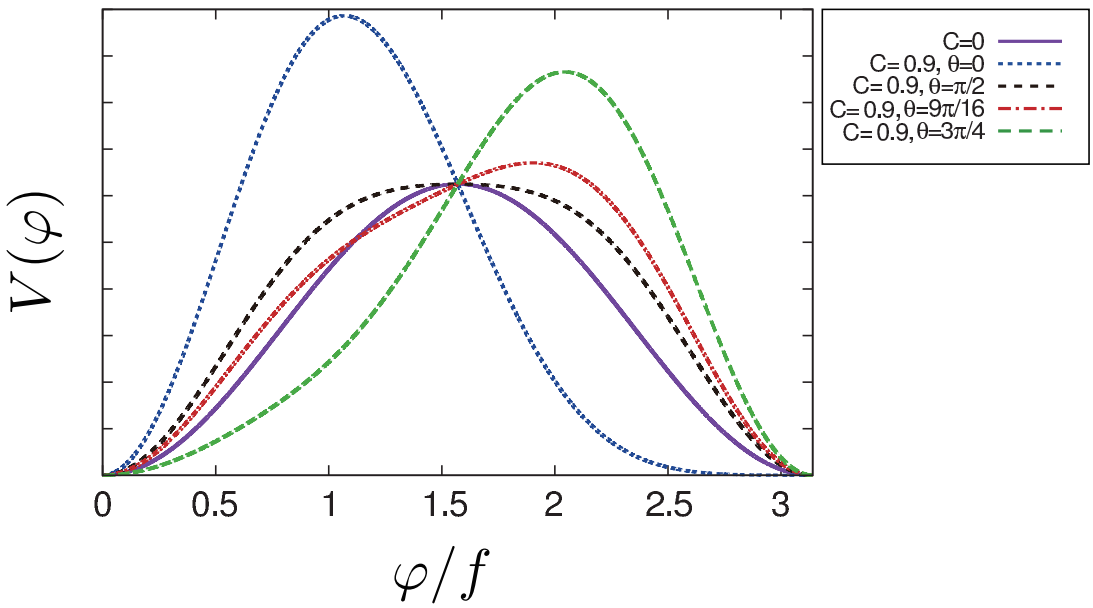

Figure 1: Shape of the scalar potential of φ for $C = 0.9$ with $\theta = 0, \pi/2, 9\pi/16$, and $3\pi/4$ for fixed f and M. The case of $C = 0$ is also shown for comparison.

and calculated the slow-roll parameters,

$$\epsilon = \frac{M_P^2}{2}\left(\frac{V'}{V}\right)^2, \quad \eta = M_P^2 \frac{V''}{V}, \tag{2.9}$$

at 60 e-folding before inflation ends. Then the scalar spectral index and tensor-to-scalar ratio are given by

$$n_s = 1 - 6\epsilon + 2\eta, \quad r = 16\epsilon. \tag{2.10}$$

The results are shown in Fig. 2 for the same parameters as in Fig. 1. For each line, we have varied f in the range $5M_P < f < 100M_P$. One can see that the predicted n_s and r significantly differ from those of natural inflation, and that they can lie in the region favored by the Planck result, $n_s = 0.9655 \pm 0.058$ and $r < 0.09$ [9]. The Planck normalization on the primordial curvature perturbation can be satisfied by adjusting the oveall scale of the inflaton potential, which results in $M = \mathcal{O}(10^{13})$ GeV unless r is smaller than $\mathcal{O}(10^{-3})$.

2.2 Stability of inflationary path

In Ref. [11] it was shown that the inflationary trajectory may be destabilized in the presence of a coupling like $\lambda X H_u H_d$ with X being a stabilizer field, because the $H_u H_d$ direction becomes tachyonic during inflation. This constrains the coupling constant as $|\lambda| \lesssim 10^{-6}$. In our case, the stabilizer field N_2 has large yukawa couplings to $L_\alpha H_u$ to reproduce the

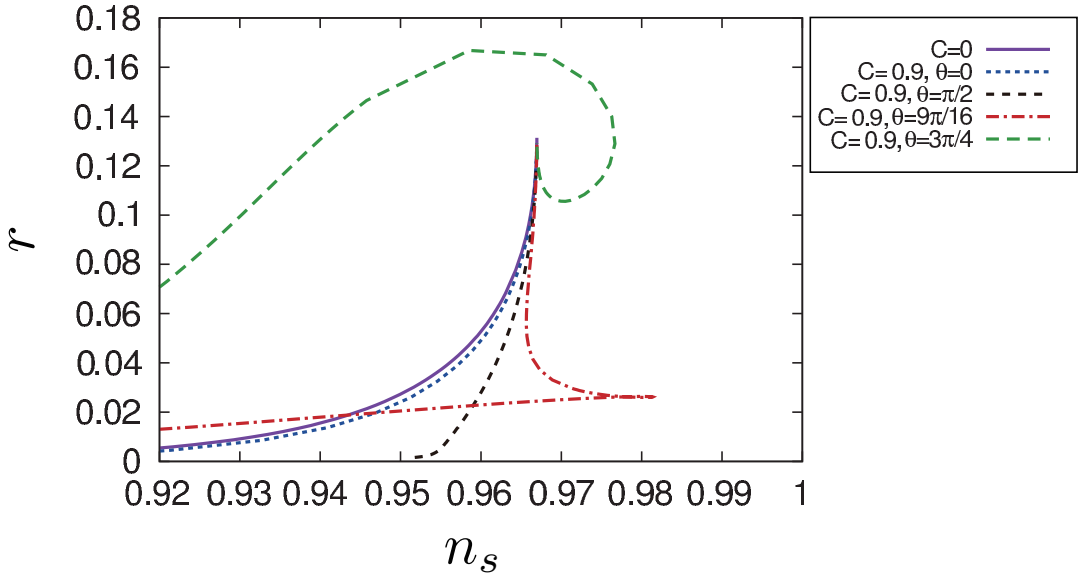

Figure 2: The prediction of (n_s, r) for the same parameters in Fig. 1. For each line, we varie f in the range $5M_P < f < 100M_P$.

observed neutrino masses, and hence one may think that this induces a similar instability. Below we show that this instability does not exist thanks to the large yukawa couplings of N_1 to $L_\alpha H_u$.

The scalar potential of L_α and H_u up to the quadratic terms during inflation is given by

$$V = \left[M N_1^* \left(\sum_\alpha y_{2\alpha} L_\alpha \right) H_u + \text{h.c.} \right] + \left| \sum_\alpha y_{1\alpha} L_\alpha \right|^2 |N_1|^2 + \sum_\alpha |y_{1\alpha}|^2 |N_1|^2 |H_u|^2, \quad (2.11)$$

where we have taken $N_2 = 0$. This is rewritten as

$$V = (M \overline{y_2} N_1^* L_2' H_u + \text{h.c.}) + \overline{y_1}^2 |N_1|^2 \left(|L_1'|^2 + |H_u|^2 \right), \quad (2.12)$$

where

$$\overline{y_1} \equiv \sqrt{\sum_\alpha |y_{1\alpha}|^2}, \quad \overline{y_2} \equiv \sqrt{\sum_\alpha |y_{2\alpha}|^2}, \quad (2.13)$$

$$L_1' \equiv \frac{1}{\overline{y_1}} \left(\sum_\alpha y_{1\alpha} L_\alpha \right), \quad L_2' \equiv \frac{1}{\overline{y_2}} \left(\sum_\alpha y_{2\alpha} L_\alpha \right). \quad (2.14)$$

Thus L_1' obtains a mass of $\overline{y_1}|N_1|$ and hence is stabilized. The mass matrix of $(H_u, L_2'^*)$ is given by

$$m^2_{HL_2'} = \begin{pmatrix} \overline{y_1}^2 |N_1|^2 & M \overline{y_2} N_1 \\ M \overline{y_2} N_1^* & kH^2 \end{pmatrix}, \quad (2.15)$$

where we have included a Hubble mass correction at the $(2,2)$ entry with k being a positive constant of order unity.[#3] This matrix has two eigenvalues of

$$M^2 \simeq \overline{y_1}^2|N_1|^2, \quad kH^2 - \frac{\overline{y_2}^2}{\overline{y_1}^2}M^2, \tag{2.16}$$

for $\overline{y_1}|N_1| \gg \overline{y_2}M$. Thus there are no tachyonic direction during inflation if $kH^2 > \overline{y_2}^2M^2/\overline{y_1}^2$. Actually this condition is easily satisfied. Although the other combination of L_α orthogonal to both L_1' and L_2' remains massless at this level, it can also have a positive Hubble mass by introducing $K = -|N_2|^2|L_\alpha|^2/M_P^2$. Therefore the inflationary path is stable and we can take $H_u = L_\alpha = 0$ during inflation.

As we shall see in Sec. 3.1, a typical value of $\overline{y}_1$ is of $\mathcal{O}(0.1)$ which may lead to super-Planckian masses for leptons and Higgs during inflation, spoiling the effective field theory description. This is our motivation to introduce a discrete shift symmetry on the inflaton field and we need $f \lesssim \mathcal{O}(10)M_P$. We note, however, that it is in principle possible to have $\overline{y}_1 \lesssim \mathcal{O}(0.01)$. In such a case, we can take $f = \mathcal{O}(100)M_P$ and the predicted (n_s, r) are close to that of the quadratic chaotic inflation.

3 Implications

3.1 Neutrino masses and mixings

Here let us show that our model can reproduce the observed neutrino masses and mixings [13, 22, 23]. To this end, it is convenient to work with a basis in which the RH neutrino masses are diagonalized:

$$W = \frac{1}{2}\tilde{M}_i\tilde{N}_i\tilde{N}_i + \tilde{y}_{i\alpha}\tilde{N}_iL_\alpha H_u, \tag{3.1}$$

with

$$\tilde{M}_1 = \tilde{M}_2 = M, \tag{3.2}$$

$$\tilde{y}_{1\alpha} = \frac{1}{\sqrt{2}}(y_{1\alpha} + y_{2\alpha}), \quad \tilde{y}_{2\alpha} = \frac{i}{\sqrt{2}}(-y_{1\alpha} + y_{2\alpha}), \tag{3.3}$$

$$\tilde{N}_1 = \frac{1}{\sqrt{2}}(N_1 + N_2), \quad \tilde{N}_2 = \frac{i}{\sqrt{2}}(N_1 - N_2), \tag{3.4}$$

where we have expanded the interactions at the potential minimum. After integrating out the RH neutrinos in (3.1), we obtain

$$W = -\frac{1}{2}\tilde{y}_{i\alpha}\tilde{y}_{j\beta}(\tilde{M}^{-1})_{ij}(L_\alpha H_u)(L_\beta H_u). \tag{3.5}$$

[#3] Other entries also receive Hubble mass corrections, but they are subdominant since $|N_1| \gtrsim M_P$ during inflation.

Thus the light neutrino mass matrix is given by

$$m^{(\nu)}_{\alpha\beta} = \frac{v^2 \sin^2\beta}{M} \tilde{y}_{i\alpha}\tilde{y}_{i\beta}, \tag{3.6}$$

where $v = 174\,\mathrm{GeV}$ and $\sin\beta \equiv \langle H_u\rangle/v$. Note that since the mass matrix $(\tilde{M}^{-1})_{ij}$ is rank 2, $m^{(\nu)}_{\alpha\beta}$ can only have two non-zero eigenvalues. Therefore, among the total 9 real parameters to characterize the neutrino mass matrix (three masses, three mixing angles and three CP phases), one mass parameter and one phase vanish. Thus we are left with 7 low energy parameters. On the other hand, $\tilde{y}_{i\alpha}$ is a general 2×3 complex matrix having 12 real parameters, but three phases can be absorbed by redefining the phase of L_α.[#4] Thus the total physical degrees of freedom in the neutrino sector is 2 (RH neutrino mass) $+(12-3)$ (neutrino yukawa: $\tilde{y}_{i\alpha}$) $= 11$ if the RH neutrino masses are taken freely. Of these, the overall rescaling $\tilde{M}_i \to \lambda_i^2 \tilde{M}_i$ combined with $\tilde{y}_{i\alpha} \to \lambda_i \tilde{y}_{i\alpha}$ does not affect the light neutrino masses. This rescaling has two parameters, hence we are left with the 9 parameters in the high energy to parameterize the light neutrino masses and mixings. This number is larger than the number of the low energy parameters, which is 7, hence a model with two RH neutrinos has enough parameters to reproduce the light neutrino mass matrix.

In our case, the RH neutrino mass matrix has only one parameter M and its value is fixed by the normalization of the primordial curvature perturbations. Since M is fixed, there are 9 parameters in the neutrino yukawa sector, which cannot be reduced further by the rescaling. Hence the situation remains intact.

By using the MNS matrix [24], neutrino mass eigenvalues are expressed as

$$m^{(\nu)}_{\bar{\alpha}}\delta_{\bar{\alpha}\bar{\delta}} = U^{(\mathrm{MNS})T}_{\bar{\alpha}\beta} m^{(\nu)}_{\beta\gamma} U^{(\mathrm{MNS})}_{\gamma\bar{\delta}}, \tag{3.7}$$

where $\bar{\alpha} = 1,2,3$ denotes the mass eigenstate basis. Here we impose $m^{(\nu)}_1 < m^{(\nu)}_2$, so that $m^{(\nu)}_1 = 0$ for the normal hierarchy (NH) and $m^{(\nu)}_3 = 0$ for the inverted hierarchy (IH). The MNS matrix is parametrized as

$$U^{(\mathrm{MNS})}_{\alpha\bar{\beta}} = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta} & c_{23}c_{13} \end{pmatrix} \times \mathrm{diag}\left(1, e^{i\alpha/2}, 1\right), \tag{3.8}$$

where $c_{ij} = \cos\theta_{ij}$, $s_{ij} = \sin\theta_{ij}$, δ is the Dirac phase and α is the Majorana phase. The neutrino yukawa is given by

$$\tilde{y}_{i\alpha}\tilde{y}_{i\beta} = \frac{M}{v^2\sin^2\beta} U^{(\mathrm{MNS})*}_{\alpha\bar{\gamma}} m^{(\nu)}_{\bar{\gamma}} \delta_{\bar{\gamma}\bar{\delta}} U^{(\mathrm{MNS})\dagger}_{\bar{\delta}\beta}. \tag{3.9}$$

[#4] The phase redefinition of L_α combined with those of right handed charged leptons $\overline{E}_\alpha$ can leave charged lepton yukawa matrix real and diagonal.

As mentioned above, this does not uniquely determine all the matrix elements of $\tilde{y}_{i\alpha}$: there are additional two degrees of freedom in the yukawa sector to determine the light neutrino mass matrix. This can be explicitly seen by solving (3.9) as [22, 25]

$$\tilde{y}_{i\alpha} = \frac{M^{1/2}}{v \sin\beta} R_{i\bar{\gamma}} \sqrt{m_{\bar{\gamma}}^{(\nu)}} \delta_{\bar{\gamma}\bar{\delta}} U_{\bar{\delta}\alpha}^{(\mathrm{MNS})\dagger}, \tag{3.10}$$

where

$$R_{i\bar{\gamma}} = \begin{pmatrix} 0 & \cos z & -\sin z \\ 0 & \sin z & \cos z \end{pmatrix} \quad \text{for} \quad \text{NH}, \tag{3.11}$$

and

$$R_{i\bar{\gamma}} = \begin{pmatrix} -\sin z & \cos z & 0 \\ \cos z & \sin z & 0 \end{pmatrix} \quad \text{for} \quad \text{IH}, \tag{3.12}$$

with z being an arbitrary complex parameter, corresponding to the additional degrees of freedom.

The best-fit values of the observed parameters are [26]

$$\Delta m_{12}^2 = 7.54 \times 10^{-5}\,\mathrm{eV}^2, \quad \Delta m_{23}^2 = 2.43 \times 10^{-3}\,\mathrm{eV}^2, \tag{3.13}$$

$$\sin^2\theta_{12} = 0.308, \quad \sin^2\theta_{23} = 0.437, \quad \sin^2\theta_{13} = 2.34 \times 10^{-2}. \tag{3.14}$$

for NH, and

$$\Delta m_{12}^2 = 7.54 \times 10^{-5}\,\mathrm{eV}^2, \quad \Delta m_{23}^2 = 2.38 \times 10^{-3}\,\mathrm{eV}^2, \tag{3.15}$$

$$\sin^2\theta_{12} = 0.308, \quad \sin^2\theta_{23} = 0.455, \quad \sin^2\theta_{13} = 2.40 \times 10^{-2}. \tag{3.16}$$

for IH. We can determine the yukawa matrix $y_{i\alpha}$ by using these values for arbitrary values of δ, α and z. Note that $\overline{y_1}$ and $\overline{y_2}$ in (2.13) are independent of δ and α. Also they are independent of z as long as z is real and in such a case we have $\overline{y_1} = \overline{y_2}$. If z is real, we obtain

$$\overline{y_1} = \overline{y_2} = 0.139 \quad \text{for} \quad \text{NH}, \tag{3.17}$$

$$\overline{y_1} = \overline{y_2} = 0.180 \quad \text{for} \quad \text{IH}, \tag{3.18}$$

for $M = 2 \times 10^{13}\,\mathrm{GeV}$ and $\sin\beta = 1$. We regard them as "typical" values. On the other hand, if z has an imaginary component, the prediction changes. For example, for $z = i$ we obtain

$$\overline{y_1} = 0.0511, \quad \overline{y_2} = 0.377 \quad \text{for} \quad \text{NH}, \tag{3.19}$$

$$\overline{y_1} = 0.0660, \quad \overline{y_2} = 0.488 \quad \text{for} \quad \text{IH}. \tag{3.20}$$

By taking a large value of imaginary component of z, we can make a hierarchy between $\overline{y_1}$ and $\overline{y_2}$. For $|z| \lesssim \mathcal{O}(1)$, $\overline{y_1}$ and $\overline{y_2}$ are of $\mathcal{O}(0.01 - 1)$.

3.2 Reheating

Now let us consider the reheating after inflation. In our model, the reheating process is slightly nontrivial because of the large yukawa coupling of the inflaton. Just after inflation, the inflaton φ begins a coherent oscillation with its amplitude of order M_P, and the coupled charged leptons and Higgs (and their superpartners) obtain masses of order $\sim y\varphi$ where y collectively denotes the neutrino yukawa coupling. Since this is much larger than the inflaton mass M, the perturbative decay of the inflaton is not kinematically allowed. Instead, non-perturbative particle production, called preheating, happens when φ passes the origin $\varphi \simeq 0$ [27]. The produced particles decay into lighter ones before the inflaton again moves back to $\varphi \simeq 0$ [28]. For example, Higgs boson decays into quarks through yukawa couplings.

The effective decay rate of φ through this process is estimated as [29]

$$\Gamma_\phi \sim \frac{2y^2 M}{\pi^{7/2} g} \equiv bM, \tag{3.21}$$

where g collectively denotes the Higgs coupling to lighter particles and $b \sim \mathcal{O}(0.01 - 0.1)$. Thus after a few Hubble time after inflation, a significant fraction of the inflaton energy density is transferred to radiation:

$$\rho_r \sim b\rho_\phi. \tag{3.22}$$

If this process continues to produce radiation even after the radiation energy density begins to dominate the universe, the reheating temperature is given by

$$T_{\rm R} \sim 10^{15}\,{\rm GeV} \left(\frac{b}{0.01}\right)^{1/2} \left(\frac{M}{2\times 10^{13}\,{\rm GeV}}\right)^{1/2}. \tag{3.23}$$

On the other hand this preheating process may become ineffective due to the thermal mass correction to the Higgs particles. Then the main process that transfers the inflaton energy to radiation becomes the scattering of light particles in thermal plasma with inflaton. In our case, at the high inflaton amplitude regime $y\varphi \gtrsim T$, it is the effective inflaton coupling with SU(2) gauge bosons after integrating out heavy Higgs and leptons that is responsible for such an effective dissipation rate. The oscillation-averaged dissipation rate is estimated as [29]

$$\Gamma_\phi^{(\rm dis)} \sim \frac{b' y \alpha_W^2 T^2}{\tilde{\varphi}}, \tag{3.24}$$

where b' is an order one numerical constant, α_W is the SU(2) fine structure constant and $\tilde{\varphi}$ denotes the oscillation amplitude of φ. The inflaton is thermalized when this rate becomes comparable to the Hubble scale. This occurs at $H \sim (\alpha_W^2 b^{1/2} b' y)^{2/3} M \sim \mathcal{O}(10^{-3})M$. Thus the reheating temperature in this case is estimated to be

$$T_{\rm R} \sim \mathcal{O}(0.01) \times \sqrt{MM_P} \sim 10^{14}\,{\rm GeV} \left(\frac{M}{2\times 10^{13}\,{\rm GeV}}\right)^{1/2}. \tag{3.25}$$

In either case, the reheating temperature is so high that the RH sneutrino inflaton becomes thermalized.

3.3 Leptogenesis

Now let us consider implication for the leptogenesis scenario [3]. In our model, two RH neutrino masses are nearly degenerate, and so, we must take account of the resonant effect [30–35]. It should be noticed that the effective CP asymmetry in the RH neutrino decay vanishes in the exact degenerate limit.[#5] Thus we need a small diagonal mass matrix element of $\sim \delta M$ ($|\delta M| \ll M$) in (2.4).[#6] After diagonalizing the RH neutrino mass matrix, we obtain

$$W = \frac{1}{2}\tilde{M}_i \tilde{N}_i \tilde{N}_i + \tilde{y}_{i\alpha} \tilde{N}_i L_\alpha H_u, \tag{3.26}$$

$$\tilde{M}_1 = M - \delta M, \quad \tilde{M}_2 = M + \delta M. \tag{3.27}$$

The lepton asymmetry is generated via the CP asymmetric decay of $\tilde{N}_1$ and $\tilde{N}_2$. The CP asymmetry parameter is given by [35]

$$\epsilon_i = \frac{\mathrm{Im}\left[(\tilde{y}_{i\alpha}\tilde{y}^\dagger_{\alpha j})^2\right]}{8\pi(\tilde{y}_{i\alpha}\tilde{y}^\dagger_{\alpha i})} \frac{\tilde{M}_i \tilde{M}_j (\tilde{M}_i^2 - \tilde{M}_j^2)}{(\tilde{M}_i^2 - \tilde{M}_j^2)^2 + (\tilde{M}_i \Gamma_i - \tilde{M}_j \Gamma_j)^2}, \tag{3.28}$$

where Γ_i denotes the decay width of $\tilde{N}_i$. The ϵ_i parameter is maximized and can be $\mathcal{O}(1)$ for $\tilde{M}_1^2 - \tilde{M}_2^2 \sim \tilde{M}_1 \Gamma_1$ assuming $\tilde{M}_1 \Gamma_1 \gtrsim \tilde{M}_2 \Gamma_2$.

The final baryon asymmetry, after the sphaleron conversion of the lepton number to the baryon number, is given by

$$\frac{n_B}{s} = \frac{8}{23}\frac{n_L}{s} = \frac{8}{23}\kappa \sum_{i=1,2} \epsilon_i \frac{n_{N_i}}{s}, \tag{3.29}$$

where κ represents the suppression factor due to the washout effect [36]. In our case, RH neutrinos are expected to be in thermal equilibrium at $T \sim M_i$: $\Gamma_{N_1}/H_{T=M_i} \sim \mathcal{O}(100)$. Thus we have $\kappa \sim \mathcal{O}(0.01)$. To explain the observed value $n_B/s \simeq 9 \times 10^{-11}$, we need $\epsilon_i \sim 10^{-5}$.[#7]

3.4 Gravitino problem

Finally we discuss the gravitino problem. In general, there are two contributions to the gravitino production: thermal production and nonthermal production. The nonthermal gravitino production rate from the direct decay of the inflaton depends on the inflaton

[#5] This can be seen in (2.4) that we can rotate a phase of N_1 and N_2 without affecting the RH mass term, while it can absorb the phase of yukawa matrix which appears in the CP asymmetric decay of RH neutrinos. This phase rotation is allowed only for the degenerate case, i.e., there are no diagonal elements in the mass matrix of RH neutrinos.

[#6] As long as $|\delta M| \ll 10^{-2} M$, it does not much affect the inflaton dynamics.

[#7] For $|\tilde{M}_1^2 - \tilde{M}_2^2| \ll |\tilde{M}_i \Gamma_i - \tilde{M}_j \Gamma_j|$, the calculation based on the Boltzmann equation becomes invalid. Instead, we may need Kadanoff-Baym approach to estimate the lepton asymmetry. In any case, we can obtain small ϵ_i in a degenerate limit [35].

vacuum expectation value [37–41]. In our model, thanks to the Z_2 symmetry, the inflaton VEV is zero and hence there is no significant production of the gravitino from the inflaton decay.

Thermal gravitino production, on the other hand, is very efficient in the present model because of high reheating temperature. The gravitino abundance in terms of the number to the entropy density ratio is estimated as [42]

$$\frac{n_{3/2}}{s} \simeq 2 \times 10^{-12} \left(1 + \frac{m_{\tilde{g}}^2}{3m_{3/2}^2}\right) \left(\frac{T_{\rm R}}{10^{10}\,{\rm GeV}}\right), \tag{3.30}$$

where $m_{\tilde{g}}$ and $m_{3/2}$ denote the gluino mass and gravitino mass, respectively. The cosmological consequences depend on the gravitino mass. For the unstable gravitino, in order to avoid the constraint from big-bang nucleosynthesis (BBN) [43], we need $m_{3/2} \sim 100\,{\rm TeV}$ so that it decays before BBN begins. However, even in such a case, the lightest SUSY particles (LSP) produced by the gravitino decay give a too large contribution to the relic dark matter abundance independently of the gravitino mass, if we assume the anomaly mediation relation between the gaugino and gravitino mass [44]. A solution to the LSP overproduction is to introduce a small R-parity violation so that LSP decays quickly before BBN. Another way to avoid the gravitino problem is to assume an ultra light gravitino scenario in which $m_{3/2} \lesssim 16\,{\rm eV}$ [45, 46].

4 Conclusions

In this paper we have revisited a chaotic inflation model in which the inflaton is identified with a RH sneutrino. We have introduced only two RH neutrinos, which is minimal in the sense that at least two heavy fields are required for both successful chaotic inflation and explaining the neutrino masses and mixings. To avoid super-Planckian masses of the leptons and Higgs during inflation, we have assumed that the shift symmetry of the inflaton is not completely broken by the superpotential interactions, but there remain an unbroken discrete shift symmetry, which modifies the shape of the inflaton potential. As a result, we can obtain the prediction of (n_s, r) within 1σ range of the Planck result if the decay constant f is $\mathcal{O}(10)M_P$.

Since we have only two RH neutrinos, one of the light neutrinos is massless. Although the rate of neutrinoless double beta decay is suppressed compared to the (quasi) degenerate case, it is within the reach of forthcoming experiments because $\langle m_{ee} \rangle$ is bounded below [47–49]. Inflaton has rather large yukawa couplings to explain observed neutrino masses and hence the reheating automatically happens. The predicted reheating temperature is high and leptogenesis works efficiently.

Acknowledgments

T. T. Y. thanks Prof. Raymond Volkas for the hospitality during his stay at the University of Melbourne. This work was supported by the Grant-in-Aid for Scientific Research on Scientific Research A (No.26247042 [KN and FT]), Scientific Research B (No. 26287039 [FT and TTY]), Young Scientists B (No.26800121 [KN], No. 24740135 [FT]) and Innovative Areas (No.26104009 [KN], No.15H05888 [KN], No. 23104008 [FT], and No.15H05889 [FT]). This work was supported by World Premier International Research Center Initiative (WPI Initiative), MEXT, Japan.

References

[1] A. D. Linde, Phys. Lett. B **129** (1983) 177.

[2] T. Yanagida, Conf. Proc. C **7902131**, 95 (1979); M. Gell-Mann, P. Ramond and R. Slansky, Conf. Proc. C **790927**, 315 (1979) [arXiv:1306.4669 [hep-th]]; P. Minkowski, Phys. Lett. B **67**, 421 (1977).

[3] M. Fukugita and T. Yanagida, Phys. Lett. B **174**, 45 (1986).

[4] H. Murayama, H. Suzuki, T. Yanagida and J. Yokoyama, Phys. Rev. Lett. **70**, 1912 (1993). Phys. Rev. D **50**, 2356 (1994) [hep-ph/9311326].

[5] K. Nakayama, F. Takahashi and T. T. Yanagida, Phys. Lett. B **730**, 24 (2014) [arXiv:1311.4253 [hep-ph]].

[6] H. Murayama, K. Nakayama, F. Takahashi and T. T. Yanagida, Phys. Lett. B **738**, 196 (2014) [arXiv:1404.3857 [hep-ph]].

[7] M. Kawasaki, M. Yamaguchi and T. Yanagida, Phys. Rev. Lett. **85**, 3572 (2000) [hep-ph/0004243].

[8] R. Kallosh and A. Linde, JCAP **1011**, 011 (2010) [arXiv:1008.3375 [hep-th]].

[9] P. A. R. Ade *et al.* [Planck Collaboration], arXiv:1502.02114 [astro-ph.CO].

[10] K. Nakayama, F. Takahashi and T. T. Yanagida, Phys. Lett. B **725**, 111 (2013) [arXiv:1303.7315 [hep-ph]].

[11] K. Nakayama, F. Takahashi and T. T. Yanagida, JCAP **1308**, 038 (2013) [arXiv:1305.5099 [hep-ph]].

[12] K. Nakayama, F. Takahashi and T. T. Yanagida, Phys. Lett. B **737**, 151 (2014) [arXiv:1407.7082 [hep-ph]].

[13] P. H. Frampton, S. L. Glashow and T. Yanagida, Phys. Lett. B **548**, 119 (2002) [hep-ph/0208157].

[14] M. Czerny and F. Takahashi, Phys. Lett. B **733**, 241 (2014) [arXiv:1401.5212 [hep-ph]].

[15] M. Czerny, T. Higaki and F. Takahashi, JHEP **1405**, 144 (2014) [arXiv:1403.0410 [hep-ph]].

[16] M. Czerny, T. Higaki and F. Takahashi, Phys. Lett. B **734**, 167 (2014) [arXiv:1403.5883 [hep-ph]].

[17] T. Kobayashi and F. Takahashi, JCAP **1101**, 026 (2011) [arXiv:1011.3988 [astro-ph.CO]].

[18] M. Czerny, T. Kobayashi and F. Takahashi, Phys. Lett. B **735**, 176 (2014) [arXiv:1403.4589 [astro-ph.CO]].

[19] T. Higaki and F. Takahashi, JHEP **1407**, 074 (2014) [arXiv:1404.6923 [hep-th]]; Phys. Lett. B **744**, 153 (2015) [arXiv:1409.8409 [hep-ph]].

[20] R. Kappl, H. P. Nilles and M. W. Winkler, arXiv:1511.05560 [hep-th].

[21] K. Choi and H. Kim, arXiv:1511.07201 [hep-th].

[22] A. Ibarra and G. G. Ross, Phys. Lett. B **591**, 285 (2004) [hep-ph/0312138].

[23] K. Harigaya, M. Ibe and T. T. Yanagida, Phys. Rev. D **86**, 013002 (2012) [arXiv:1205.2198 [hep-ph]].

[24] Z. Maki, M. Nakagawa and S. Sakata, Prog. Theor. Phys. **28**, 870 (1962).

[25] J. A. Casas and A. Ibarra, Nucl. Phys. B **618**, 171 (2001) [hep-ph/0103065].

[26] F. Capozzi, G. L. Fogli, E. Lisi, A. Marrone, D. Montanino and A. Palazzo, Phys. Rev. D **89**, 093018 (2014) [arXiv:1312.2878 [hep-ph]].

[27] L. Kofman, A. D. Linde and A. A. Starobinsky, Phys. Rev. Lett. **73**, 3195 (1994) [hep-th/9405187]; Phys. Rev. D **56**, 3258 (1997) [hep-ph/9704452].

[28] G. N. Felder, L. Kofman and A. D. Linde, Phys. Rev. D **59**, 123523 (1999) [hep-ph/9812289].

[29] K. Mukaida and K. Nakayama, JCAP **1301**, 017 (2013) [arXiv:1208.3399 [hep-ph]]; JCAP **1303**, 002 (2013) [arXiv:1212.4985 [hep-ph]]; K. Mukaida, K. Nakayama and M. Takimoto, JHEP **1312**, 053 (2013) [arXiv:1308.4394 [hep-ph]].

[30] L. Covi and E. Roulet, Phys. Lett. B **399**, 113 (1997) [hep-ph/9611425].

[31] A. Pilaftsis, Nucl. Phys. B **504**, 61 (1997) [hep-ph/9702393].

[32] W. Buchmuller and M. Plumacher, Phys. Lett. B **431**, 354 (1998) [hep-ph/9710460].

[33] A. Pilaftsis and T. E. J. Underwood, Nucl. Phys. B **692**, 303 (2004) [hep-ph/0309342]; Phys. Rev. D **72**, 113001 (2005) [hep-ph/0506107].

[34] A. Anisimov, A. Broncano and M. Plumacher, Nucl. Phys. B **737**, 176 (2006) [hep-ph/0511248].

[35] M. Garny, A. Kartavtsev and A. Hohenegger, Annals Phys. **328**, 26 (2013) [arXiv:1112.6428 [hep-ph]].

[36] G. F. Giudice, A. Notari, M. Raidal, A. Riotto and A. Strumia, Nucl. Phys. B **685**, 89 (2004) [hep-ph/0310123].

[37] M. Kawasaki, F. Takahashi and T. T. Yanagida, Phys. Lett. B **638**, 8 (2006) [hep-ph/0603265]; Phys. Rev. D **74**, 043519 (2006) [hep-ph/0605297].

[38] T. Asaka, S. Nakamura and M. Yamaguchi, Phys. Rev. D **74**, 023520 (2006) [hep-ph/0604132].

[39] M. Endo, K. Hamaguchi and F. Takahashi, Phys. Rev. D **74**, 023531 (2006) [hep-ph/0605091].

[40] M. Endo, F. Takahashi and T. T. Yanagida, Phys. Lett. B **658**, 236 (2008) [hep-ph/0701042].

[41] M. Endo, F. Takahashi and T. T. Yanagida, Phys. Rev. D **76**, 083509 (2007) [arXiv:0706.0986 [hep-ph]].

[42] M. Bolz, A. Brandenburg and W. Buchmuller, Nucl. Phys. B **606**, 518 (2001) [Erratum-ibid. B **790**, 336 (2008)] [hep-ph/0012052]; J. Pradler and F. D. Steffen, Phys. Rev. D **75**, 023509 (2007) [hep-ph/0608344]; Phys. Lett. B **648**, 224 (2007) [hep-ph/0612291]; V. S. Rychkov and A. Strumia, Phys. Rev. D **75**, 075011 (2007) [hep-ph/0701104].

[43] M. Kawasaki, K. Kohri and T. Moroi, Phys. Rev. D **71**, 083502 (2005) [astro-ph/0408426].

[44] L. Randall and R. Sundrum, Nucl. Phys. B **557**, 79 (1999) [hep-th/9810155]; G. F. Giudice, M. A. Luty, H. Murayama and R. Rattazzi, JHEP **9812**, 027 (1998) [hep-ph/9810442].

[45] E. Pierpaoli, S. Borgani, A. Masiero and M. Yamaguchi, Phys. Rev. D **57**, 2089 (1998) [astro-ph/9709047].

[46] M. Viel, J. Lesgourgues, M. G. Haehnelt, S. Matarrese and A. Riotto, Phys. Rev. D **71**, 063534 (2005) [astro-ph/0501562].

[47] J. J. Gomez-Cadenas, J. Martin-Albo, M. Sorel, P. Ferrario, F. Monrabal, J. Munoz-Vidal, P. Novella and A. Poves, JCAP **1106**, 007 (2011) [arXiv:1010.5112 [hep-ex]].

[48] J. Barea, J. Kotila and F. Iachello, Phys. Rev. C **87**, no. 1, 014315 (2013) [arXiv:1301.4203 [nucl-th]].

[49] G. Benato, Eur. Phys. J. C **75**, no. 11, 563 (2015) [arXiv:1510.01089 [hep-ph]].