Leptogenesis in Inflaton Decay

T. Asaka\textsuperscript{1}, K. Hamaguchi\textsuperscript{1}, M. Kawasaki\textsuperscript{2} and T. Yanagida\textsuperscript{1,2}

\textsuperscript{1}Department of Physics, University of Tokyo, Tokyo 113-0033, Japan
\textsuperscript{2}Research Center for the Early Universe, University of Tokyo, Tokyo 113-0033, Japan

(June 15, 1999)

Abstract

We study a leptogenesis via decays of heavy Majorana neutrinos produced non-thermally in inflaton decays. We find that this scenario is fully consistent with existing supersymmetric inflation models such as for topological or for hybrid inflation and the Froggatt-Nielsen mechanism generating hierarchies in quark and lepton mass matrices. The reheating temperature $T_R$ of inflation may be taken as low as $T_R \simeq 10^8$ GeV to avoid the cosmological gravitino problem.
1 Introduction

Primordial lepton asymmetry is converted to baryon asymmetry \[1\] in the early universe through the “sphaleron” effects of the electroweak gauge theory \[2\] if it is produced before the electroweak phase transition. Therefore, lepton-number violation at high energies may be an important ingredient for creating the baryon (matter-antimatter) asymmetry in the present universe. A generic low-energy prediction of the leptogenesis is the presence of effective operators for the lepton-number violation such as

\[
\mathcal{O} = \frac{f_{ij}}{M} l_i l_j H H, \tag{1}
\]

where \(M\) is the scale of lepton-number violation, \(l_i (i=1,2,3)\) lepton doublets and \(H\) the Higgs scalar field in the standard model.

The above operators induce small neutrino masses \[3\] in the true vacuum \(\langle H \rangle \simeq 246/\sqrt{2} \text{ GeV}\). There is now a convincing experimental evidence \[4\] that neutrinos have indeed small masses \(m_\nu\) of order 0.01–0.1 eV. Thus, the leptogenesis scenario \[1\] seems to be the most plausible mechanism for creating the cosmological baryon asymmetry.

The lepton-number violating operators Eq. (1) arise from the exchange of heavy Majorana neutrinos \(N_i (i=1,2,3)\). Decays of these heavy neutrinos produce very naturally lepton asymmetry if \(C\) and \(CP\) are not conserved \[1\]. There have been, so far, proposed and investigated various scenarios for the leptogenesis depending on production mechanisms of the heavy Majorana neutrinos \(N_i \) \[1, 5, 6, 7\].

In this letter we discuss a leptogenesis scenario where \(N_i\) are produced non-thermally in inflaton decays \[3\]. We find that this scenario is fully consistent with existing inflation models such as for topological \[3\] or for hybrid inflation \[3\], and the Froggatt-Nielsen (FN) mechanism \[10\] generating hierarchies in the quark and lepton mass matrices. The reheating temperature \(T_R\) of inflation may be taken as low as \(T_R \simeq 10^8 \text{ GeV}\) avoiding the cosmological gravitino problem \[11\].

We assume supersymmetry (SUSY) throughout this letter. In SUSY models there is an interesting leptogenesis mechanism \[12\] of the Affleck-Dine type whose detailed analysis will be given elsewhere.
2 Lepton asymmetry in the decays of heavy Majorana neutrinos

Decays of the heavy neutrinos $N$ into leptons $l$ and Higgs $H_u$ doublets violate lepton-number conservation since they possess two decay channels:

$$N \rightarrow H_u + l$$
$$N \rightarrow \overline{H_u} + \overline{l}. \quad (2)$$

We consider only the $N_1$ decay, provided the mass $M_1$ of $N_1$ is smaller than the others ($M_1 \ll M_2, M_3$). Interference between amplitudes of the tree-level and one-loop diagrams yields a lepton asymmetry $\epsilon_1$ in Eq. (3) by using an effective $CP$-violating phase $\delta_{\text{eff}}$ as

$$\epsilon_1 = \frac{\Gamma(N_1 \rightarrow H_u + l) - \Gamma(N_1 \rightarrow \overline{H_u} + \overline{l})}{\Gamma(N_1 \rightarrow H_u + l) + \Gamma(N_1 \rightarrow \overline{H_u} + \overline{l})}$$
$$= -\frac{3}{16\pi \left(h_\nu h_\nu^\dagger\right)_{11}} \left[ \text{Im} \left(h_\nu h_\nu^\dagger\right)_{13}^2 M_1 + \text{Im} \left(h_\nu h_\nu^\dagger\right)_{12}^2 M_3 \right]. \quad (4)$$

Here we have taken a basis where the mass matrix for $N_i$ is diagonal and the Yukawa coupling constants $(h_\nu)_{ij}$ are defined in the superpotential as $W = (h_\nu)_{ij} N_i l_j H_u$. We have included both of one-loop vertex and self-energy corrections [14].

The FN model considered in the next section suggests $\left(h_\nu h_\nu^\dagger\right)_{12}^2 / M_3 \simeq \left(h_\nu h_\nu^\dagger\right)_{13}^2 / M_2$, and hence we rewrite the lepton asymmetry parameter $\epsilon_1$ in Eq. (4) by using an effective $CP$-violating phase $\delta_{\text{eff}}$ as

$$\epsilon_1 \simeq \frac{3\delta_{\text{eff}}}{16\pi \left(h_\nu h_\nu^\dagger\right)_{11}} \left| \left(h_\nu h_\nu^\dagger\right)_{13}^2 \right| \frac{M_1}{M_3}. \quad (5)$$

Assuming $|(h_\nu)_{i3}| > |(h_\nu)_{i2}| \gg |(h_\nu)_{i1}|$ ($i = 1, 3$) we obtain

$$\epsilon_1 \simeq \frac{3\delta_{\text{eff}}}{16\pi} \left| (h_\nu)_{33}^2 \right| \frac{M_1}{M_3}$$
$$\simeq \frac{3\delta_{\text{eff}} m_{\nu_3} M_1}{16\pi \langle H_u \rangle^2}. \quad (6)$$

Here, we have used the see-saw formula [3] [see Eq. (13)]

$$m_{\nu_3} \simeq \frac{\left| (h_\nu)_{33}^2 \right| \langle H_u \rangle^2}{M_3}. \quad (7)$$

1 Here $N, l(\overline{l})$ and $H_u(\overline{H_u})$ denote fermionic or bosonic components of corresponding supermultiplets.
Taking the maximal $CP$ violating phase $|\delta_{\text{eff}}| \simeq 1$, $m_{\nu_3} \simeq 3 \times 10^{-2}$ eV suggested\footnote{The atmospheric neutrino oscillation observed in the Superkamiokande experiments indicates the neutrino mass squared difference $m_{\nu_3}^2 - m_{\nu_2}^2 \simeq (0.5-6) \times 10^{-3}$ eV\footnote{In the SUSY standard model $\langle H_u \rangle = \sin \beta \langle H \rangle$. We assume, here, tan $\beta \equiv \langle H_u \rangle / \langle H_d \rangle \simeq O(1)$, where $H_u$ and $H_d$ are Higgs supermultiplets couple to up-type and down-type quarks, respectively. If one takes tan $\beta \simeq 50$ the lepton asymmetry is reduced only by factor 2.}} from the Superkamiokande experiments \cite{4}, and $\langle H_u \rangle \simeq 174/\sqrt{2}$ GeV\footnote{The atmospheric neutrino oscillation observed in the Superkamiokande experiments indicates the neutrino mass squared difference $m_{\nu_3}^2 - m_{\nu_2}^2 \simeq (0.5-6) \times 10^{-3}$ eV\footnote{In the SUSY standard model $\langle H_u \rangle = \sin \beta \langle H \rangle$. We assume, here, tan $\beta \equiv \langle H_u \rangle / \langle H_d \rangle \simeq O(1)$, where $H_u$ and $H_d$ are Higgs supermultiplets couple to up-type and down-type quarks, respectively. If one takes tan $\beta \simeq 50$ the lepton asymmetry is reduced only by factor 2.}} we get \cite{3}

$$\epsilon_1 \simeq -10^{-6} \left( \frac{M_1}{10^{10} \text{ GeV}} \right). \quad (8)$$

Let us now consider the production of heavy neutrino $N_1$ through inflaton $\varphi$ decays, which leads to a constraint on the inflaton mass $m_\varphi$ as

$$m_\varphi > 2M_1. \quad (9)$$

We consider the case where $M_1 \gtrsim 100 T_R$ ($T_R$ is the reheating temperature of the inflation)\footnote{The atmospheric neutrino oscillation observed in the Superkamiokande experiments indicates the neutrino mass squared difference $m_{\nu_3}^2 - m_{\nu_2}^2 \simeq (0.5-6) \times 10^{-3}$ eV\footnote{In the SUSY standard model $\langle H_u \rangle = \sin \beta \langle H \rangle$. We assume, here, tan $\beta \equiv \langle H_u \rangle / \langle H_d \rangle \simeq O(1)$, where $H_u$ and $H_d$ are Higgs supermultiplets couple to up-type and down-type quarks, respectively. If one takes tan $\beta \simeq 50$ the lepton asymmetry is reduced only by factor 2.}} In this case the Majorana neutrino $N_1$ is always out of thermal equilibrium even if it has $O(1)$ Yukawa coupling to $H_u$ and $l_i$, and the $N_1$ behaves like frozen-out, relativistic particle with the energy $E_{N_1} \simeq m_\varphi/2$. As we will see in the next section, the $N_1$ decays immediately after produced by the inflaton decays and hence we obtain lepton-to-entropy ratio \cite{6}

$$\frac{n_L}{s} \simeq \frac{3}{2} \epsilon_1 B_r \frac{T_R}{m_\varphi} \simeq -10^{-6} B_r \left( \frac{T_R}{10^{10} \text{ GeV}} \right) \left( \frac{M_1}{m_\varphi} \right), \quad (10)$$

where $B_r$ is the branching ratio of the inflaton decay into $N_1$ channel. We impose $T_R \lesssim 10^8$ GeV which is required to suppress sufficiently the density of the gravitino \cite{11}. Notice that the reheating temperature $T_R$ is bounded from below, $T_R \gtrsim 10^6$ GeV, otherwise the produced lepton asymmetry is too small as $n_L/s < 10^{-10}$.

The lepton asymmetry in Eq. \cite{10} is converted to the baryon asymmetry through the “sphaleron” effects which is given by

$$\frac{n_B}{s} \simeq a \frac{n_L}{s}, \quad (11)$$

with $a \simeq -8/23$ \cite{10} in the minimal SUSY standard model. To explain the observed baryon asymmetry

$$\frac{n_B}{s} \simeq (0.1-1) \times 10^{-10}, \quad (12)$$

\cite{2}
The Froggatt-Nielsen (FN) mechanism [10] is the most attractive framework for explaining the observed hierarchies in the quark and charged lepton mass matrices, which is based on a broken $U(1)_F$ symmetry. A gauge singlet field $\Phi$ carrying the FN charge $Q_\Phi = -1$ is assumed to have a vacuum-expectation value $\langle \Phi \rangle$ and then the Yukawa couplings of Higgs supermultiplets arise from nonrenormalizable interactions of $\Phi$ as

$$W = g_{ij} \left( \frac{\Phi}{M_G} \right)^{Q_i + Q_j} \Psi_i \Psi_j H_{u(d)},$$

where $Q_i$ are the $U(1)_F$ charges of various supermultiplets $\Psi_i$, $g_{ij}$ $\mathcal{O}(1)$ coupling constants, and the gravitational scale $M_G \simeq 2.4 \times 10^{18}$ GeV. The mass hierarchies for quarks and charged leptons are well explained in terms of their FN charges with $\epsilon \equiv \langle \Phi \rangle / M_G \simeq 1/17$ [17].

We adopt the above mechanism to the neutrino sector. Possible FN charges for various supermultiplets are shown in Table 1. The charges for $e_c^i$ are taken to be the same as those of up-type quarks assuming that they belong to the same 10’s in the $SU(5)$ grand unified theory. The charge $a$ of the $l_3$ may be 0 or 1 (see Ref. [17] for details). The assignment of the FN charges for the lepton doublets yields a mass matrix for neutrinos [17, 15] as

$$\left( m_\nu \right)_{ij} \simeq \epsilon^{2a} \begin{pmatrix} \epsilon^2 & \epsilon & \epsilon \\ \epsilon & 1 & 1 \\ \epsilon & 1 & 1 \end{pmatrix} \frac{\langle H_u \rangle^2}{M_0}.$$

Here, Majorana masses $M_i$ of the heavy neutrinos $N_i$ are given by $M_i = \epsilon^{2Q_i} M_0$ with $Q_i$ being the FN charges of $N_i$. This mass matrix leads to a large $\nu_\mu - \nu_\tau$ mixing angle, which is consistent
with the atmospheric neutrino oscillation \cite{[17]}. Using \( m_{\nu_3} \simeq 3 \times 10^{-2} \text{eV} \) we may derive

\[
M_0 \simeq \epsilon^{2a} \times 10^{15} \text{GeV} \simeq \begin{cases} 
10^{15} \text{GeV} & \text{for } a = 0 \\
10^{13} \text{GeV} & \text{for } a = 1 
\end{cases}.
\]

(16)

Hereafter, we assume \( a = 0 \) and \( M_0 \simeq 10^{15} \text{GeV} \) for simplicity. Notice that the final result depends only on \( B_r, T_R \) and \( M_1/m_\phi \) as seen in Eq. (10), and hence the FN charge \( a \) itself is irrelevant to the present analysis. For \( d = 1 \) and \( d = 2 \) we obtain \( M_1 \simeq 10^{13} \text{GeV} \) and \( 10^{10} \text{GeV} \), respectively.\(^5\) The partial decay rate \( \Gamma(N_1 \to H_u l) \) is given by

\[
\Gamma(N_1 \to H_u l) \simeq \frac{1}{8\pi} |\epsilon|^{2(a+d)} M_1.
\]

(17)

For the case of \( M_1 \simeq 10^{10} \text{GeV} \) we see from Eq. (13) that the inflaton mass should lie in a range of \( 10^{10} \text{GeV} \lesssim m_\varphi \lesssim 10^{12} \text{GeV} \) to explain the present observed baryon asymmetry. We will find in the following section that a topological inflation model \cite{[8]} proposed recently satisfies this condition and \( T_R \simeq 10^8 \text{GeV} \). The total decay rate of the \( N_1 \) is \( \Gamma_{N_1} \simeq (1/4\pi) |\epsilon|^4 M_1 \simeq 10^4 \text{GeV} \) which is much larger than the inflaton decay rate \( \Gamma_\varphi \sim T_R^2/M_G \simeq 10^{-2} \text{GeV} \), and hence the \( N_1 \) decays immediately after produced by the inflaton decay.\(^6\)

For the case of \( M_1 \simeq 10^{13} \text{GeV} \) we find from Eq. (13) that the inflaton mass of \( 10^{13} \text{GeV} \lesssim m_\varphi \lesssim 10^{15} \text{GeV} \) is required, which nicely fits in a SUSY hybrid inflation model \cite{[9]} as seen in the next section. We find that \( \Gamma_{N_1} \gg \Gamma_\varphi \) in this case and our approximation used in deriving Eq. (10) is also justified.

4 Leptogenesis in inflaton decays

We are now at the point to discuss the leptogenesis in inflaton decays. We first consider the case of \( M_1 \simeq 10^{10} \text{GeV} \) and a topological inflation model with the inflaton mass of \( 10^{10} \text{GeV} \lesssim m_\varphi \lesssim 10^{12} \text{GeV} \). We adopt a SUSY topological inflation which has been proposed in Ref. \cite{[8]}. The model has the following superpotential \( W \) and Kähler potential \( K \);

\[
W = v^2 \chi(1 - g\phi^2/M_G^2),
\]

(18)

\[
K = |\chi|^2 + |\phi|^2 + \frac{1}{M_G^2} \left( k_1 |\chi|^2 |\phi|^2 - \frac{k_2}{4} |\chi|^4 \right) + \cdots,
\]

(19)

where \( v \) is the energy scale of the inflation, \( g, k_1 \) and \( k_2 \) are constants of order unity, and the ellipsis denotes higher order terms. Here, we impose \( U(1)_R \times Z_2 \) symmetry and omit higher-order terms for simplicity. We assume \( \chi(\theta) \to e^{-2i\alpha} \chi(\theta e^{i\alpha}) \) and \( \phi(\theta) \to \phi(\theta e^{i\alpha}) \) under the

\(^5\) For the case of \( a = 1 \) we obtain \( M_1 \simeq 10^{10} \text{GeV} \) for \( d = 1 \).

\(^6\) Perturbative calculation in the present analysis is justified, since \( \Gamma_{N_1} \ll M_1 \).
$U(1)_R$, and $\chi$ is even and $\phi$ is odd under the $Z_2$. This discrete $Z_2$ symmetry is an essential ingredient for the topological inflation [18].

The potential of scalar components of the supermultiplets $\chi$ and $\phi$ is obtained in the standard manner, which yields SUSY vacua

$$\langle \chi \rangle = 0, \quad \langle \phi \rangle = \pm \frac{M_G}{\sqrt{g}} \equiv \pm \eta,$$

in which the potential energy vanishes. Here, the scalar components of the supermultiplets are denoted by the same symbols as the corresponding supermultiplets.

A topological inflation [18] occurs if the vacuum-expectation value $\langle \phi \rangle$ is of order of the gravitational scale $M_G$. The critical value $\eta_{cr}$ of $\langle \phi \rangle$ for which the topological inflation occurs was investigated in Refs. [19, 20]. We adopt the result in Ref. [19], which gives

$$\eta_{cr} \simeq 1.7M_G.$$  \hspace{1cm} (21)

Thus, for topological inflation to take place, $\eta$ should be larger than $\eta_{cr}$.

The potential for the region $|\chi|, |\phi| \ll M_G$ is written approximately as

$$V \simeq v^4 [1 - g\phi^2/M_G^2]^2 + (1 - k_1)v^4|\phi|^2 + k_2v^4|\chi|^2/M_G^2.$$  \hspace{1cm} (22)

Since the $\chi$ field quickly settles down to the origin for the case $k_2 \gtrsim 1$, we set $\chi = 0$ in Eq. (22) assuming $k_2 \gtrsim 1$. For $g > 0$ and $k_1 < 1$, we identify the inflaton field $\varphi/\sqrt{2}$ with the real part of the field $\phi$ since the imaginary part of $\phi$ has a positive mass larger than the size of the negative mass of $\varphi$ near the origin $\phi \simeq 0$. Then, we obtain a potential for the inflaton as

$$V(\varphi) \simeq v^4 - \frac{k}{2M_G^2}v^4\varphi^2,$$  \hspace{1cm} (23)

where $k \equiv 2g + k_1 - 1$. The inflaton has a mass $m_\varphi$

$$m_\varphi \simeq 2|\sqrt{g}v^2/M_G| = 2v^2/\eta,$$  \hspace{1cm} (24)

in the true vacuum Eq. (20).

The slow-roll condition for the inflation is satisfied for $0 < \kappa < 1$ and $\varphi < \varphi_f$ where $\varphi_f$ is of order of $M_G$, which provides the value of $\varphi$ at the end of inflation. The scale factor of the universe increases by a factor of $e^N$ when the inflaton $\varphi$ rolls slowly down the potential from $\varphi_N$ to $\varphi_f$. The $e$-fold number $N$ is given by

$$N \simeq \int_{\varphi_f}^{\varphi_N} d\varphi \frac{V}{V'}M_G^2 \simeq \frac{1}{\kappa} \ln \frac{\varphi_f}{\varphi_N}.$$  \hspace{1cm} (25)

6
The amplitude of primordial density fluctuations $\delta \rho / \rho$ due to the inflation is written as

$$\frac{\delta \rho}{\rho} \simeq \frac{1}{5 \sqrt{3 \pi M_G^3}} \frac{V^{3/2}(\varphi_N)}{|V'(|\varphi_N|)|} \simeq \frac{1}{5 \sqrt{3 \pi} \kappa \varphi_N M_G} v^2.$$  

This should be normalized to the data on anisotropies of the cosmic microwave background radiation (CMBR) by the Cosmic Background Explorer (COBE) satellite [21]. Since the $e$-fold number $N$ corresponding to the COBE scale is about 60, the COBE normalization gives

$$\frac{V^{3/2}(\varphi_{60})}{M_G^3 |V'(|\varphi_{60}|)|} \simeq 5.3 \times 10^{-4}.$$  

In this model the spectrum index $n_s$ is given by (for details see Ref. [8])

$$n_s \simeq 1 - 2\kappa.$$  

Since the COBE data implies $n_s$ as $n_s = 1.0 \pm 0.2$ [21], we should take $\kappa \lesssim 0.1$. Along with Eqs. (25) and (27), we obtain

$$v \simeq 2.3 \times 10^{-2} \sqrt{\kappa \varphi_f M_G} e^{-\frac{\kappa N}{2}},$$

which leads to

$$m_\varphi \simeq 3.8 \times 10^{11} \text{GeV},$$

for $\kappa \simeq 0.1$, $\eta \simeq \eta_{cr}$ and $\varphi_f \simeq M_G$. [7]

The inflaton decay into $2N_1$ occurs through nonrenormalizable interactions in the Kähler potential such as

$$K = \sum_i C_i |\phi|^2 |\psi_i|^2 / M_G^2,$$

where $\psi_i$ denote supermultiplets for SUSY standard-model particles including $N_1$, and $C_i$ are coupling constants of order unity. With these interactions the decay rate $\Gamma_\varphi$ of the inflaton is estimated as

$$\Gamma_\varphi \simeq \sum_i \frac{C_i^2 \eta^2 m_\varphi^3}{8 \pi M_G^4},$$

which yields the reheating temperature $T_R$ given by

$$T_R \simeq 0.092 C \eta (m_\varphi / M_G)^{3/2} \simeq 10^8 \text{GeV},$$

for $C = \sqrt{\sum_{i} |C_i|^2} \simeq 4.3$, $\kappa \simeq 0.1$, and $\eta \simeq \eta_{cr}$. Therefore, this topological inflation model can provide $T_R \simeq 10^8$ GeV and $m_\varphi \simeq 10^{11-12}$ GeV, which leads to

$$\frac{n_L}{s} \simeq -10^{-8} B_r \left( \frac{M_1}{m_\varphi} \right) \simeq -10^{-10},$$

If one takes $\kappa \lesssim 0.07$, one gets $m_\phi \gtrsim 10^{12}$ GeV.
for $M_1 \simeq 10^{10}$ GeV and $B_r \simeq 0.1$–$1$.

Next, we study a hybrid inflation model \cite{9} which gives the inflaton mass between $10^{13}$ GeV and $10^{15}$ GeV and $T_R \simeq 10^8$ GeV. The SUSY hybrid inflation model contains two kinds of supermultiplets: one is $\phi$ and the others are $\Psi$ and $\overline{\Psi}$. The model is also based on the $U(1)_R$ symmetry. The superpotential and Kähler potential are given by \cite{23, 24, 9}

$$W = -\mu^2 \phi + \lambda \phi \overline{\Psi} \Psi,$$  \hspace{1cm} (35)

$$K = |\phi|^2 + |\Psi|^2 + |\overline{\Psi}|^2 + \cdots,$$  \hspace{1cm} (36)

where the ellipsis denotes higher-order terms, which we neglect in the present analysis. To satisfy the $D$-term flatness condition we take always $\Psi = \overline{\Psi}$ in our analysis assuming an extra $U(1)$ gauge symmetry.\cite{9} As shown in Ref.\cite{9} the real part of $\phi$ is identified with the inflaton field $\varphi/\sqrt{2}$. The potential is minimized at $\Psi = \overline{\Psi} = 0$ when $\varphi$ is larger than $\varphi_c = \sqrt{2}\mu/\sqrt{\lambda}$, and inflation occurs for $\varphi_c < \varphi < M_G$.

Including one-loop corrections \cite{23}, the potential for the inflaton $\varphi$ is given by

$$V \simeq \mu^4 \left[ 1 + \frac{\lambda^2}{8\pi^2} \ln \left( \frac{\varphi}{\varphi_c} \right) \right].$$  \hspace{1cm} (37)

The $e$-fold number $N$ is estimated as

$$N \simeq \frac{4\pi^2}{\lambda^2 M_G^2} \left( \varphi_N^2 - \varphi_c^2 \right).$$  \hspace{1cm} (38)

For $\lambda \gtrsim 10^{-3}$ the scales $\mu$ and $\langle \Psi \rangle \equiv \xi$ are determined by the COBE normalization [Eq. (27)] as

$$\mu \simeq 5.5 \times 10^{15} \lambda^{1/2} \text{GeV},$$  \hspace{1cm} (39)

$$\xi \simeq 5.5 \times 10^{15} \text{GeV},$$  \hspace{1cm} (40)

where we have used $N \simeq 60$ at the COBE scale.\cite{9} Then, the mass of the inflaton $m_\varphi$ in the true minimum ($\langle \Psi \rangle = \langle \overline{\Psi} \rangle = \xi$ and $\langle \phi \rangle = 0$) is given by

$$m_\varphi \simeq \sqrt{2}\mu \simeq 7.7 \times 10^{15} \lambda \text{GeV},$$  \hspace{1cm} (41)

\footnote{Leptogenesis in the hybrid inflation model has been considered in Ref. \cite{22}, where the lepton asymmetry arises from the decays of the heavy neutrino $N_2$ in the second family.}

\footnote{One may consider the following superpotential instead of Eq. (35); $W = -\mu^2 \phi + \lambda \phi \Psi^2 + \lambda' \phi \Psi^3 / M_G + \cdots$. In this case one does not need the extra $U(1)$ gauge symmetry. The third term is introduced to erase unwanted domain walls.}

\footnote{The spectrum index is estimated as $n_s \simeq 1$.}
which gives \( m_\varphi \simeq 10^{13} - 10^{15} \) GeV for \( \lambda \simeq 10^{-3} - 10^{-1} \). We assume that the inflaton decays through nonrenormalizable interactions in the Kähler potential as

\[
K = \sum_i C'_i |\Sigma|^2 |\psi_i|^2 / M_G^2,
\]

with \( \Sigma \equiv (\Psi + \overline{\Psi})/\sqrt{2} \). The reheating temperature is estimated as

\[
T_R \simeq 0.092 C' \xi (m_\varphi / M_G)^{3/2} \simeq 10^8 \text{GeV},
\]

for \( C' = \sqrt{\sum_i |C_i|^2} \simeq 40 - 0.04 \). Notice that \( m_\Sigma = m_\varphi \) in the true vacuum. The obtained reheating temperature \( T_R \simeq 10^8 \) GeV and inflaton mass \( m_\varphi \simeq 10^{13} - 10^{15} \) GeV lead to the required lepton asymmetry, \( (n_L / s) \simeq -10^{-8} B_r (M_1 / m_\varphi) \simeq -10^{-10} \), for \( M_1 \simeq 10^{13} \) GeV and \( B_r \simeq 0.01 - 1 \).

We should comment on an interesting possibility that the extra \( U(1) \) gauge symmetry introduced to fix \( \Psi = \overline{\Psi} \) is nothing but the \( B - L \) symmetry, which is spontaneously broken by the condensations \( \langle \Psi \rangle = \langle \overline{\Psi} \rangle = \xi \)

Namely, the Majorana masses \( M_i \) of the heavy neutrinos \( N_i \) are induced by Yukawa couplings

\[
W = g_i N_i N_i \Psi.
\]

In this model the inflaton \( \varphi \) and \( \Sigma \) decay mainly through the Yukawa couplings Eq. (44) and the reheating temperature is given by

\[
T_R \simeq 0.092 g_1 \sqrt{m_\Sigma M_G},
\]

provided \( m_\varphi = m_\Sigma < 2 M_{2,3} \). We see that the \( T_R \) is too high as \( T_R \gtrsim 10^{12} \) GeV for \( M_1 \simeq 2 g_1 \langle \Psi \rangle \simeq 10^{13} \) GeV. Thus, we consider \( M_1 \simeq 10^{10} \) GeV. In this case we get \( T_R \simeq 10^8 \) GeV for \( m_\varphi \simeq 10^{12} \) GeV \(^{12}\) which yields the desired lepton asymmetry \( (n_L / s) \simeq -10^{-10} \), since \( B_r \simeq 1 \).

The small inflaton mass \( m_\varphi \) or equivalently the small coupling \( \lambda \simeq 10^{-4} \) in Eq. (35) may be naturally explained by a large FN charge of \( \phi \) (e.g. \( Q_\phi = 3 \)).

\(^{11}\) The spontaneous breakdown of the \( B - L \) gauge symmetry produces cosmic strings. Astrophysical analyses on anisotropies of the CMBR give a constraint \( \xi \lesssim 6.3 \times 10^{15} \) GeV \(^{24}\). Thus, the present parameter region Eq. (40) may be accessible to future satellite experiments \(^{26}\).

\(^{12}\) The \( B - L \) charges of \( N_i, \Psi \) and \( \overline{\Psi} \) are \(-1, +2\) and \(-2\), respectively. In Ref. \(^{22}\) the \( B - L \) charges of \( \Psi \) and \( \overline{\Psi} \) are chosen as \(+1\) and \( -1 \) so that one has nonrenormalizable interactions \( W = g_i N_i N_i \Psi / M_G \).

\(^{13}\) If we take \( m_\varphi \simeq 10^{10} \) GeV (\( > 2 M_1 \)), the reheating temperature \( T_R \) becomes lower as \( T_R \simeq 10^6 - 10^7 \) GeV. The desired lepton asymmetry \( (n_L / s) \simeq -10^{-10} \) may be obtained even in this case, because of \( m_\varphi / M_1 \simeq 2 \) and \( B_r \simeq 1 \) [see Eq. (10)].
5 Conclusions

We have studied the SUSY leptogenesis via the decay of the heavy Majorana neutrino $N_1$’s which are produced non-thermally in inflaton decays. Generic argument has shown that the leptogenesis works without overproduction of gravitinos if the reheating temperature $T_R \simeq 10^8$ GeV and the inflaton mass $m_\phi$ satisfies $2M_1 < m_\phi \lesssim 100M_1$. With use of the Froggatt-Nielsen model the mass of the heavy Majorana neutrino $N_1$ in the first family is easily chosen as $M_1 = 10^{10}$ GeV or $M_1 = 10^{13}$ GeV depending on the $U(1)_F$ charge for $N_1$. We have found that the required inflaton masses and reheating temperature are naturally realized in existing SUSY inflation models such as for topological [8] or for hybrid inflation [9].

Acknowledgements

The author would like to thank T. Kanazawa for useful discussion on hybrid inflation models. This work is partially supported by the Japan Society for the Promotion of Science (TA,KH) and “Priority Area: Supersymmetry and Unified Theory of Elementary Particles (♯707)” (MK,TY).
References

[1] M. Fukugita and T. Yanagida, Phys. Lett. B174 (1986) 45.

[2] V.A. Kuzmin, V.A. Rubakov, and M.E. Shaposhnikov, Phys. Lett. B155 (1985) 36.

[3] T. Yanagida, in Proc. Workshop on the unified theory and the baryon number in the universe, (Tsukuba, 1979), eds. O. Sawada and S. Sugamoto, Report KEK-79-18 (1979); M. Gell-Mann, P. Ramond and R. Slansky, in “Supergravity” (North-Holland, Amsterdam, 1979) eds. D.Z. Freedman and P. van Nieuwenhuizen.

[4] Y. Fukuda et al. [Superkamiokande Collaboration], Phys. Lett. B436 (1998) 33; Phys. Rev. Lett. 81 (1998) 1562.

[5] See, for a recent review, W. Buchmüller and M. Plüümacher, hep-ph/9904310; M. Plüümacher, Nucl. Phys. B530 (1998) 207 and reference therein.

[6] B.A. Campbell, S. Davidson and K.A. Olive, Nucl. Phys. B399 (1993) 111; K. Kumekawa, T. Moroi and T. Yanagida, Prog. Theor. Phys. 92 (1994) 437; G. Lazarides, hep-ph/9904428 and reference therein; G.F. Giudice, M. Peloso, A. Riotto, I. Tkachev, hep-ph/9905242.

[7] H. Murayama and T. Yanagida, Phys. Lett. B322 (1994) 349; H. Murayama, H. Suzuki, T. Yanagida and J. Yokoyama, Phys. Rev. Lett. 70 (1993) 1912, Phys. Rev. D50 (1994) 2356.

[8] K.I. Izawa, M. Kawasaki and T. Yanagida, Prog. Theor. Phys. 101 (1999) 1129.

[9] C. Panagiotakopoulos, Phys. Rev. D55 (1997) 7335; A. Linde and A. Riotto, Phys. Rev. D56 (1997) 1841.

[10] C.D. Froggatt and H.B. Nielsen, Nucl. Phys. B147 (1979) 277.

[11] M.Y. Khlopov and A.D. Linde, Phys. Lett. B138 (1984) 265; J. Ellis, J.E. Kim and D.V. Nanopoulos, Phys. Lett. B145 (1984) 181; M. Kawasaki and T. Moroi, Prog. Theor. Phys. 93 (1995) 879; see also, for example, a recent analysis, E. Holtmann, M. Kawasaki, K. Kohri and T. Moroi, Phys. Rev. D60 (1999) 023506.

[12] H. Murayama and T. Yanagida, in Ref. [7].

[13] L. Covi, E. Roulet and F. Vissani, Phys. Lett. B384 (1996) 169; M. Flanz, E.A. Paschos and U. Sarkar, Phys. Lett. B345 (1995) 248; Phys. Lett. B384 (1996) 487 (E).

[14] W. Buchmüller and M. Plüümacher, Phys. Lett. B431 (1998) 354.

[15] W. Buchmüller and T. Yanagida, Phys. Lett. B445 (1999) 399.

[16] S.Y. Khlebnikov and M.E. Shaposhnikov, Nucl. Phys. B308 (1988) 885; J.A. Harvey and M.S. Turner, Phys. Rev. D42 (1990) 3344.
[17] J. Sato and T. Yanagida, Talk given at 18th International Conference on Neutrino Physics and Astrophysics (NEUTRINO 98), (Takayama, Japan, 1998), hep-ph/9809307; P. Raymond, Talk given at 18th International Conference on Neutrino Physics and Astrophysics (NEUTRINO 98), (Takayama, Japan, 1998), hep-ph/9809401.

[18] A. Linde, Phys. Lett. B327 (1994) 208; A. Vilenkin, Phys. Rev. Lett. 72 (1994) 3137.

[19] N. Sakai, H. Shinkai, T. Tachizawa and K. Maeda, Phys. Rev. D53 (1996) 655; (E) Phys. Rev. D54 (1996) 2981.

[20] I. Cho and A. Vilenkin, Phys. Rev. D56 (1998) 7621; A.A. de Laix, M. Trodden and T. Vachaspati Phys. Rev. D57 (1998) 7186.

[21] C.L. Bennett et al., Astrophys. J. 464 (1996) L1.

[22] G. Lazarides, in Ref. [5].

[23] G. Dvali, Q. Shafi and R.K. Shaefer, Phys. Rev. Lett. 73 (1994) 1886.

[24] E.J. Copeland, A.R. Liddle, D.H. Lyth, E.D. Stewart and D. Wands, Phys. Rev. D49 (1994) 6410.

[25] B. Allen, R.R. Caldwell, S. Dodelson, L. Knox, E.P.S. Shellard, and A. Stebbins, Phys. Rev. Lett. 79 (1997) 2624.

[26] http://map.gsfc.nasa.gov/; http://astro.estec.esa.nl/SA-general/Projects/Planck.