Leptogenesis from an $\tilde{N}$-dominated early universe

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We investigate in detail leptogenesis by the decay of a coherent right-handed sneutrino $\tilde{N}$ having dominated the energy density of the early universe, which was originally proposed by two of the authors (H.M. and T.Y.). Once the $\tilde{N}$ dominant universe is realized, the amount of generated lepton asymmetry (and hence baryon asymmetry) is determined only by the properties of the right-handed neutrino, regardless of the history before it dominates the universe. Moreover, thanks to the entropy production by the decay of the right-handed sneutrino, thermally produced relics are sufficiently diluted. In particular, the cosmological gravitino problem can be avoided even when the reheating temperature of inflation is higher than 10^{10} GeV, in a wide range of the gravitino mass $m_{3/2} \simeq 10$ MeV-100 TeV. If the gravitino mass is in the range $m_{3/2} \simeq 10$ MeV-1 GeV as in some gauge-mediated supersymmetry breaking models, the dark matter in our Universe can be dominantly composed of the gravitino. The quantum fluctuation of $\tilde{N}$ during inflation causes an isocurvature fluctuation which may be detectable in the future.

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I. INTRODUCTION

Neutrino oscillations, especially the atmospheric neutrino oscillation observed in the SuperKamiokande experiments [1], is one of the greatest discoveries in the field of particle physics after the success of the standard model. The data suggest small but finite masses of neutrinos. Such small neutrino masses can be naturally obtained via the seesaw mechanism [2], implying the existence of lepton number violation. There has been, therefore, growing interest in leptogenesis [3] as a production mechanism of the baryon asymmetry in the present universe. In fact, the "sphaleron" process [4] converts lepton asymmetry into baryon asymmetry, and nonzero lepton asymmetry can be produced by the decay of the heavy right-handed neutrino [3].

On the other hand, supersymmetry (SUSY) has been regarded as an attractive candidate for physics beyond the standard model, since it protects the huge hierarchy between the electroweak and unification scales against radiative corrections as well as leads to a beautiful unification of the gauge coupling constants. In Ref. [5], H.M. and T.Y. proposed new possibilities for leptogenesis in the framework of SUSY. Under the assumption of SUSY, there appears a very simple and attractive mechanism to produce the lepton asymmetry,[1] that is, the condensation of the scalar component of the right-handed neutrino and its decay into the leptons and antileptons.

In this paper, we investigate in detail the leptogenesis by the decay of a coherent right-handed sneutrino. In particular, we discuss the case in which the coherent oscillation of the right-handed sneutrino dominates the energy density of the early universe. It is extremely interesting that the amount of produced baryon asymmetry is determined mainly by the decay rate of the right-handed neutrino, whatever happened before the coherent oscillation dominates the universe. Furthermore, as a big bonus, thermally produced gravitinos are diluted by the entropy production due to the decay of the coherent right-handed sneutrino, so that the cosmological gravitino problem [9–12] can be avoided even when the reheating temperature T_R of the inflation is higher than 10^{10} GeV, in a wide range of the gravitino mass $m_{3/2} \sim 10$ MeV-100 TeV.

In particular, this dilution of the thermally produced gravitinos has great advantages in the gauge-mediated SUSY breaking (GMSB) models [13]. The GMSB mechanism has been regarded as a very attractive candidate for the SUSY breaking, since it suppresses quite naturally the flavor changing processes, which are inherent problems in the SUSY standard model. In general, GMSB models predict that the gravitino is the lightest SUSY particle[2] and

[1]Another interesting possibility for leptogenesis with SUSY proposed in Ref. [5] is the leptogenesis via the flat direction including the charged lepton doublet L [5–7], which is based on the Affleck-Dine mechanism [8].

[2]This is not the case if the SUSY breaking is mediated by a bulk gauge field in higher-dimensional space-time [14].

stable.[3] Usually, the relic abundance of the gravitino is proportional to the reheating temperature, and there are severe upper bounds on the reheating temperature T_R depending on the gravitino mass $m_{3/2}$, in order to avoid that the energy density of the gravitino overclose the present universe [12].[4] In our scenario, however, this overclosure bound is completely removed because of the aforementioned right-handed sneutrino decay, and a reheating temperature even higher than 10^{10} GeV is possible for $m_{3/2} \gtrsim 10$ MeV. Furthermore, as we will see, the present energy density of the gravitino is determined independently of the reheating temperature, and the gravitino mass can be predicted as $m_{3/2} \simeq 10$ MeV-1 GeV from the baryon asymmetry in the present universe, if the dominant component of the dark matter is the gravitino.

II. LEPTOGENESIS BY COHERENT RIGHT-HANDED SNEUTRINO

A. The MSSM with right-handed neutrinos

Let us start by introducing three generations of heavy right-handed neutrinos N_i with masses M_i to the minimal supersymmetric standard model (MSSM), which have a superpotential

$$W = \frac{1}{2} M_i N_i N_i + h_{i\alpha} N_i L_\alpha H_u, \tag{1}$$

where L_α ($\alpha = e, \mu, \tau$) and H_u denote the supermultiplets of the lepton doublets and the Higgs doublet which couples to up-type quarks, respectively. The small neutrino mass is obtained by integrating out the heavy right-handed neutrinos, which is given by [2]

$$(m_\nu)_{\alpha\beta} = -\sum_i h_{i\alpha} h_{i\beta} \frac{\langle H_u \rangle^2}{M_i}. \tag{2}$$

During inflation, the scalar component of the right-handed neutrino $\tilde{N}$ can acquire a large amplitude [5,15,16] if the Hubble expansion rate of the inflation $H_{\rm inf}$ is larger than the mass of the $\tilde{N}$. Let us assume that there exists (at least) one right-handed neutrino with a mass lighter than $H_{\rm inf}$, and that it develops a large expectation value during the inflation. Hereafter, we focus on the lightest right-handed sneutrino $\tilde{N}_1$ for simplicity. (Possible contributions from the heavier right-handed sneutrinos, $\tilde{N}_2$ and $\tilde{N}_3$ will be discussed at the end of this section.) It is assumed here that the potential for the right-handed neutrino is given simply by the mass term

$$V = M_1^2 |\tilde{N}_1|^2 \tag{3}$$

and L and H_u vanish.[5]

After the end of the inflation, the Hubble parameter H decreases with cosmic time t as $H \propto t^{-1}$, and $\tilde{N}_1$ begins to oscillate around the origin when H becomes smaller than the mass of the right-handed sneutrino M_1. Then, the coherent oscillation eventually decays when $H = \Gamma_{N_1}$ ($t \sim \Gamma_{N_1}^{-1}$), where $\Gamma_{N_1} = (1/4\pi) \Sigma_\alpha |h_{1\alpha}|^2 M_1$ is the decay rate of the $\tilde{N}_1$. Because $\tilde{N}_1$ decays into leptons (and Higgs) as well as their antiparticles, its decay can produce lepton-number asymmetry if CP is not conserved [3]. The generated lepton number density is given by

$$n_L = \epsilon_1 M_1 |\tilde{N}_{1d}|^2, \tag{4}$$

where $|\tilde{N}_{1d}|$ is the amplitude of the oscillation when it decays, and ϵ_1 denotes the lepton-asymmetry parameter in the decay of $\tilde{N}_1$. Assuming a mass hierarchy $M_1 \ll M_2, M_3$ in the right-handed neutrino sector, the explicit form of ϵ_1 is given by [17]

$$\begin{aligned} \epsilon_1 &\equiv \frac{\Gamma(\tilde{N}_1 \to L + H_u) - \Gamma(\tilde{N}_1 \to \overline{L} + \overline{H_u})}{\Gamma(\tilde{N}_1 \to L + H_u) + \Gamma(\tilde{N}_1 \to \overline{L} + \overline{H_u})} \\ &\simeq -\frac{3}{8\pi} \frac{1}{(hh^\dagger)_{11}} \sum_{i=2,3} \mathrm{Im}[(hh^\dagger)_{1i}^2] \frac{M_1}{M_i}. \end{aligned} \tag{5}$$

Here, L and H_u ($\overline{L}$ and $\overline{H_u}$) symbolically denote fermionic or scalar components of corresponding supermultiplets (and their antiparticles). By using the seesaw formula in Eq. (2), this ϵ_1 parameter can be rewritten in terms of the heaviest neutrino mass m_{ν_3} and an effective CP violating phase $\delta_{\rm eff}$ [18]:

$$\begin{aligned} \epsilon_1 &= \frac{3}{8\pi} \frac{M_1}{\langle H_u \rangle^2} \frac{\mathrm{Im}[h(m_\nu^*)h^T]_{11}}{(hh^\dagger)_{11}} \\ &\equiv \frac{3}{8\pi} \frac{M_1}{\langle H_u \rangle^2} m_{\nu_3} \delta_{\rm eff} \\ &\simeq 2 \times 10^{-10} \left(\frac{M_1}{10^6 \ \mathrm{GeV}} \right) \left(\frac{m_{\nu_3}}{0.05 \ \mathrm{eV}} \right) \delta_{\rm eff}. \end{aligned} \tag{6}$$

Here, we have used $\langle H_u \rangle = 174$ GeV$\times \sin\beta$, where $\tan\beta \equiv \langle H_u \rangle / \langle H_d \rangle$. ($H_d$ is the Higgs field which couples to down-type quarks.) Here and hereafter, we take $\sin\beta \simeq 1$ for simplicity. As for the heaviest neutrino mass, we take $m_{\nu_3} \simeq 0.05$ eV as a typical value, suggested from the atmospheric neutrino oscillation observed in the SuperKamiokande experiments [1].

B. Cosmic lepton asymmetry

The fate of the generated lepton asymmetry depends on whether or not the coherent oscillation of $\tilde{N}_1$ dominates the

[3]We assume here that the R parity is exact.

[4]For a very light gravitino $m_{3/2} \lesssim 1$ keV, there is no gravitino problem [11].

[5]The parameters we prefer (as we will see later) give a large effective mass to L and H_u because $\tilde{N}_1 \sim M_{pl}$. Therefore, vanishing L and H_u is natural.

energy density of the universe before it decays [5]. In this paper, we mainly discuss the leptogenesis scenario from the universe dominated by $\tilde{N}_1$. (We will give a brief comment on the case where $\tilde{N}_1$ does not dominate the universe in the Appendix.) As we shall show soon, once the $\tilde{N}_1$ dominant universe is realized, the present baryon asymmetry is determined only by the properties of the right-handed neutrino, whatever happened before the $\tilde{N}_1$ dominates the universe. We first derive the amount of the generated lepton asymmetry just assuming that the $\tilde{N}_1$ dominates the universe, and after that we will discuss the necessary conditions of the present scenario.

Once $\tilde{N}_1$ dominates the universe before it decays, the universe is reheated again at $H=\Gamma_{N_1}$ by the decay of $\tilde{N}_1$. The energy density of the resulting radiation, with a temperature T_{N_1}, is given by the following relation:

$$\frac{\pi^2}{30} g_* T_{N_1}^4 = M_1^2 |\tilde{N}_{1d}|^2 = 3 M_{pl}^2 \Gamma_{N_1}^2, \tag{7}$$

while the entropy density is given by

$$s = \frac{2\pi^2}{45} g_* T_{N_1}^3. \tag{8}$$

Here, $M_{pl}=2.4\times 10^{18}$ GeV is the reduced Planck scale and g_* is the number of effective degrees of freedom, which is $g_* \simeq 200$ for temperatures $T \gg 1$ TeV in the SUSY standard model. From the above equations, the ratio of the lepton number density to the entropy density is given by the following simple form:

$$\frac{n_L}{s} = \frac{3}{4}\epsilon_1 \frac{T_{N_1}}{M_1} \simeq 1.5\times 10^{-10} \left(\frac{T_{N_1}}{10^6\ \text{GeV}}\right)\left(\frac{m_{\nu_3}}{0.05\ \text{eV}}\right)\delta_{\text{eff}}. \tag{9}$$

Here, we have required that the decay of the $\tilde{N}_1$ occurs in an out-of-equilibrium way, namely, $T_{N_1} < M_1$, so that the produced lepton-number asymmetry not be washed out by lepton-number violating interactions mediated by N_1.

Because the lepton asymmetry is produced before the electroweak phase transition at $T\sim 100$ GeV, it is partially converted [3] into the baryon asymmetry through the ''sphaleron'' effects [4]

$$\frac{n_B}{s} = a\,\frac{n_L}{s}, \tag{10}$$

where $a = -\frac{8}{23}$ in the SUSY standard model [19]. This ratio takes a constant value as long as an extra entropy production does not take place at a later epoch. Therefore, as mentioned in the Introduction, the baryon asymmetry in the present universe is indeed determined only by the decay temperature of the right-handed sneutrino T_{N_1} (and the effective CP violating phase δ_{eff}), given in Eq. (9). Thus it is independent of unknown parameters of the inflation such as the reheating temperature T_R. Assuming the effective CP violating phase $\delta_{\text{eff}}\,(\leqslant 1)$ to be not too small, the observed baryon asymmetry $n_B/s \simeq (0.4-1)\times 10^{-10}$ [20] is obtained by taking

$$T_{N_1} \simeq 10^6 - 10^7\ \text{GeV}. \tag{11}$$

Now let us recall the conditions we have required so far. We have required the following two conditions. (i) $\tilde{N}_1$ dominates the universe before it decays and (ii) $\tilde{N}_1$ decays in an out-of-equilibrium way. By taking the T_{N_1} in Eq. (11), the condition of the out-of-equilibrium decay is given by

$$M_1 > T_{N_1} \simeq 10^6 - 10^7\ \text{GeV}. \tag{12}$$

Notice that the temperature T_{N_1} is determined by the decay rate of the $\tilde{N}_1$ [see Eq. (7)], and hence is related to the mass and couplings of $\tilde{N}_1$. The relation is given by

$$\sqrt{\sum_\alpha |h_{1\alpha}|^2} \simeq 5\times 10^{-6} \left(\frac{T_{N_1}}{10^6\ \text{GeV}}\right)^{1/2} \left(\frac{T_{N_1}}{M_1}\right)^{1/2}. \tag{13}$$

Thus, we need Yukawa couplings $h_{1\alpha}$ which are as small as the electron Yukawa coupling.

C. Conditions for $\tilde{N}$ dominance

In order to discuss whether or not $\tilde{N}_1$ dominates the universe, it is necessary to consider the history of the universe before it decays. Here, we assume that the potential of the $\tilde{N}_1$ is ''flat'' up to the Planck scale, namely, the potential is just given by the mass term $M_1^2|\tilde{N}_1|^2$ up to the Planck scale. (This may not be the case when the masses of the right-handed neutrinos are induced by a breaking of an additional gauge symmetry. We will discuss such a case in the next section.)

Assuming the flatness of the $\tilde{N}_1$'s potential up to the Planck scale (i.e., only the mass term), the initial amplitude of the oscillation is naturally given by $|\tilde{N}_{1i}| \simeq M_{pl}$, since above the Planck scale the scalar potential is expected to be exponentially lifted by the supergravity effects.[6] Then, the energy density of $\tilde{N}_1$ when it starts the coherent oscillation is given by $\rho_{N_1} \simeq M_1^2 M_{pl}^2$.

The rest of the total energy density of the universe at $H = M_1$ is dominated by (i) the oscillating inflation ψ or (ii) the radiation, depending on the decay rate of the inflaton Γ_ψ. If $\Gamma_\psi < M_1$, the reheating process of the inflaton has not com-

[6] Even though it is possible that $\tilde{N}_1$ has a larger initial amplitude $|\tilde{N}_{1i}| > M_{pl}$ (see, e.g., Ref. [21]), it depends on the scalar potential beyond the Planck scale, so that we do not discuss this possibility in this paper.

pleted yet at $H = M_1$, and the inflaton ψ is still oscillating around its minimum, whose energy density is given by $\rho_\psi \simeq 2 M_1^2 M_{pl}^2$. The ratio of the energy density of $\tilde{N}_1$ to that of the inflaton, $\rho_{N_1}/\rho_\psi \simeq 1/2$, takes a constant value until either of these oscillations decays. Because the energy density of the radiation $\rho_{\rm rad}$ resulting from the inflaton decay is diluted faster than ρ_{N_1}, the oscillating $\tilde{N}_1$ dominates the universe if its decay rate Γ_{N_1} is slow enough compared with that of the inflaton Γ_ψ, $\Gamma_{N_1} \ll \Gamma_\psi$.

On the other hand, if $\Gamma_\psi > M_1$, the inflaton decay has already completed before $H = M_1$, and the energy density of the radiation at $H = M_1$ is given by $\rho_{\rm rad} \simeq 2 M_1^2 M_{pl}^2$. In this case, the oscillating $\tilde{N}_1$ dominates the universe soon after it starts the oscillation and hence before its decay.[7] Therefore, the condition for $\tilde{N}_1$ to dominate the universe is just given by $\Gamma_{N_1} \ll \Gamma_\psi$. In terms of the reheating temperature T_R, it is

$$T_R \gg T_{N_1} \simeq 10^6 - 10^7 \text{ GeV}, \tag{14}$$

which is easily satisfied in various SUSY inflation models [22]. Thus, the present leptogenesis scenario from $\tilde{N}_1$ dominated early universe is almost automatic as long as the right-handed neutrino has suitable mass and couplings given in Eqs. (12) and (13).

D. Gravitino problem ameliorated

Now let us turn to consider the cosmological gravitino problem [9–12]. There are two cases; unstable and stable gravitino. When the gravitino is not the lightest SUSY particle, it has a very long lifetime, and its decay during or after the big-bang nucleosynthesis (BBN) epoch ($t \sim 1-100$ sec) might spoil the success of the- BBN. Since the abundance of the thermally produced gravitinos at reheating epoch is proportional to the reheating temperature T_R, usually there are upper bounds on the T_R depending on the gravitino mass. The bound is given by $T_R \lesssim 10^7 - 10^9$ GeV for $m_{3/2} \simeq 100$ GeV–1 TeV [9], and $B_h T_R \lesssim 10^7 - 10^9$ GeV for $m_{3/2} \simeq$ (a few $-$ 100) TeV [10], where B_h denotes branching ratio of the gravitino decay into hadrons. However, in the present scenario, the gravitino abundance is diluted by the entropy production due to the right-handed sneutrino decay. The dilution factor is given by

$$\Delta = \begin{cases} \dfrac{T_R}{2 T_{N_1}} & (\text{for } T_R < T_{R_C}), \\[2ex] \dfrac{T_{R_C}}{2 T_{N_1}} & (\text{for } T_R > T_{R_C}), \end{cases} \tag{15}$$

where

$$T_{R_C} \equiv 7 \times 10^{11} \left(\frac{M_1}{10^6 \text{ GeV}} \right)^{1/2} \text{ GeV}. \tag{16}$$

Here, $T_R < T_{R_C}$ ($T_R > T_{R_C}$) corresponds to $\Gamma_\psi < M_1$ ($\Gamma_\psi > M_1$). Thanks to this entropy production by the $\tilde{N}_1$'s decay, the constraint from the gravitino problem applies not to the reheating temperature T_R, but to an effective temperature given by

$$\begin{aligned} T_{R\,\rm eff} &\equiv \frac{1}{\Delta} T_R \\ &= \begin{cases} 2 T_{N_1}, \\ 2 T_{N_1} \left(\dfrac{T_R}{T_{R_C}} \right), \end{cases} \\ &\simeq \begin{cases} 2 \times 10^6 - 2 \times 10^7 \text{ GeV} & (\text{for } T_R < T_{R_C}), \\ 2 \times 10^6 - 2 \times 10^7 \text{ GeV} \times \left(\dfrac{T_R}{T_{R_C}} \right) & (\text{for } T_R > T_{R_C}), \end{cases} \end{aligned} \tag{17}$$

which is much below the original reheating temperature T_R. Therefore, the cosmological gravitino problem can be avoided in a wide range of the gravitino mass $m_{3/2} \simeq 100$ GeV$-$100 TeV, even if the reheating temperature T_R of the inflation is higher than 10^{10} GeV. The fact that such high reheating temperature is allowed makes it very easy to construct realistic SUSY inflation models.

On the other hand, if the gravitino is the lightest SUSY particle, as in the GMSB scenario, it is completely stable. If there is no extra entropy production after the inflation, the relic abundance of the gravitinos which are produced thermally after the inflation is given by [12]

$$\Omega_{3/2} h^2 |_{\text{without } \tilde{N}_1 \text{ decay}} \simeq 0.8 \times \left(\frac{M_3}{1 \text{ TeV}} \right)^2 \left(\frac{m_{3/2}}{10 \text{ MeV}} \right)^{-1} \left(\frac{T_R}{10^6 \text{ GeV}} \right). \tag{18}$$

Here, M_3 is the gluino mass, h is the present Hubble parameter in units of 100 km sec^{-1} Mpc^{-1} and $\Omega_{3/2} = \rho_{3/2}/\rho_c$. ($\rho_{3/2}$ and ρ_c are the present energy density of the gravitino and the critical energy density of the present universe, respectively.) It is found from Eq. (18) that the overclosure limit $\Omega_{3/2} < 1$ puts a severe upper bound on the reheating temperature T_R, depending on the gravitino mass $m_{3/2}$. However, in our scenario, the "reheating" by the coherent $\tilde{N}_1$ takes place and the relic abundance of the gravitino is obtained by dividing the original abundance in Eq. (18) by the dilution factor Δ:

$$\begin{aligned} \Omega_{3/2} h^2 |_{\text{with } \tilde{N}_1 \text{ decay}} &\simeq \frac{1}{\Delta} \Omega_{3/2} h^2 |_{\text{without } \tilde{N}_1 \text{ decay}} \\ &\simeq 0.8 \times \left(\frac{M_3}{1 \text{ TeV}} \right)^2 \left(\frac{m_{3/2}}{10 \text{ MeV}} \right)^{-1} \\ &\quad \times \left(\frac{T_{R\,\rm eff}}{10^6 \text{ GeV}} \right). \end{aligned} \tag{19}$$

[7] This is the case as long as $\Gamma_{N_1} \ll M_1$.

Therefore, again, the overclosure problem can be avoided almost independently of the reheating temperature T_R, and a reheating temperature even higher than 10^{10} GeV is possible for $m_{3/2} \lesssim 10$ MeV. Moreover, it is found from this equation that the present energy density of the gravitino is independent of the reheating temperature, in a very wide range of $T_{N_1} < T_R < T_{R_C}$. Thus, we can predict the gravitino mass by requiring that the gravitino is the dominant component of the dark matter:

$$m_{3/2} \simeq 50\ \mathrm{MeV} \times \left(\frac{M_3}{1\ \mathrm{TeV}}\right)^2 \left(\frac{\Omega_{\mathrm{matter}} h^2}{0.15}\right)^{-1} \left(\frac{T_{R\ \mathrm{eff}}}{10^6\ \mathrm{GeV}}\right)$$
$$\simeq 100\ \mathrm{MeV} - 1\ \mathrm{GeV} \times \left(\frac{M_3}{1\ \mathrm{TeV}}\right)^2 \left(\frac{\Omega_{\mathrm{matter}} h^2}{0.15}\right)^{-1}, \tag{20}$$

for $T_{N_1} < T_R < T_{R_C}$.[8] Here, we take the present matter density $\Omega_{\mathrm{matter}} \simeq 0.3$ and $h \simeq 0.7$ [25]. Notice that this prediction comes from the fact that the present energy density of the gravitino is determined by the effective temperature $T_{R\ \mathrm{eff}} = 2T_{N_1}$ (for $T_R < T_{R_C}$), while the decay temperature of the right-handed neutrino T_{N_1} is fixed by the baryon asymmetry in the present universe [see Eq. (11)].

E. Some discussions

Before closing this section, several comments are in order. The first one is about the neutrino mass m_ν. The contribution to the neutrino mass matrix $(m_\nu)_{\alpha\beta}$ from N_1 is given by

$$|(m_\nu)_{\alpha\beta}^{\mathrm{from}\ N_1}| = |h_{1\alpha} h_{1\beta}| \frac{\langle H_u \rangle^2}{M_1}$$
$$\leqslant \sum_\alpha |h_{1\alpha}|^2 \frac{\langle H_u \rangle^2}{M_1}$$
$$\simeq 7 \times 10^{-4}\ \mathrm{eV} \left(\frac{T_{N_1}}{M_1}\right)^2. \tag{21}$$

Here, we have used the relation in Eq. (13). Therefore, it is understood that the mass scale of the neutrinos suggested from the atmospheric and solar neutrino oscillations, $m_\nu \sim 10^{-1} - 10^{-3}$ eV, should be induced from the heavier right-handed neutrinos N_2 and N_3. The relative hierarchy between the mass and couplings of N_1 and those of the N_2 and N_3 might be naturally explained by a broken flavor symmetry.

For example, a broken discrete Z_6 symmetry [26] with a breaking parameter $\varepsilon \simeq \frac{1}{17}$ and charges $Q(L_\tau, L_\mu, L_e) = (a, a, a+1)$ and $Q(N_3, N_2, N_1) = (b, c, 3+d)$ gives rise to the following superpotential:

$$W = \frac{1}{2} \sum_{(i,j) \neq (1,1)} g_{ij} M_0 \varepsilon^{Q_i + Q_j} N_i N_j + \frac{1}{2} g_{11} M_0 \varepsilon^{2d} N_1 N_1 + \tilde{h}_{i\alpha} \varepsilon^{Q_i + Q_\alpha} N_i L_\alpha H_u,$$

where g_{ij} and $\tilde{h}_{i\alpha}$ are $\mathcal{O}(1)$ couplings. The above charge assignments for lepton doublets naturally lead to the realistic neutrino mass matrix including the maximal mixing for the atmospheric neutrino oscillation [27]. The overall mass scale of the right-handed neutrino M_0 is determined by $m_{\nu_3} \sim \varepsilon^{2a} \langle H_\mu \rangle^2 / M_0$. By taking $a+d=2$, this model gives $M_1 \sim \varepsilon^{2d} M_0 \sim 7 \times 10^9$ GeV, $\sqrt{\Sigma_\alpha |h_{1\alpha}|^2} \sim \varepsilon^5 \sim 7 \times 10^{-7}$, and hence $T_{N_1} \sim 1 \times 10^7$ GeV.

So far, we have considered the leptogenesis from the lightest right-handed sneutrino, $\tilde{N}_1$. The heavier right-handed sneutrino $\tilde{N}_{2(3)}$ can also develop a large amplitude during the inflation (if $M_{2(3)} < H_{\mathrm{inf}}$) and it may produce lepton asymmetry in a similar way to the $\tilde{N}_1$. However, the decay temperatures of the $\tilde{N}_2$ and $\tilde{N}_3$ cannot satisfy the out-of-equilibrium condition $T_{2(3)} < M_1$, since N_2 and N_3 must explain the mass scales of the neutrino oscillations [see Eq. (21)]. Therefore, even if the $\tilde{N}_{2(3)}$'s decay would produce additional lepton asymmetry, it would be washed out and hence it can not contribute to the resultant total lepton asymmetry.

Finally, we comment on the effects of the thermal plasma [15,28,29], which might cause an early oscillation of the right-handed sneutrino $\tilde{N}_1$ before $H = M_1$. [Notice that there is a dilute plasma with a temperature $T \simeq (T_R^2 M_{pl} H)^{1/4}$ even before the reheating process of the inflation completes [20].] There are basically two possible thermal effects. First, when the temperature T is higher than the effective mass for L and H_u, $T > m_{\mathrm{eff}} = \sqrt{\Sigma_\alpha |h_{1\alpha}|^2} |\tilde{N}_1|$, the $\tilde{N}_1$ receives an additional thermal mass $\delta M_{\mathrm{th}}^2 = (1/4) \Sigma_\alpha |h_{1\alpha}|^2 T^2$ from the Yukawa coupling to L and H_u [28]. Thus, the $\tilde{N}_1$ field would start an early oscillation if the additional thermal mass becomes larger than the Hubble expansion rate before $H = M_1$. However, even if $\tilde{N}_1$ receives the thermal mass, the ratio of the thermal mass to the Hubble expansion rate is given by

$$\frac{\delta M_{\mathrm{th}}^2}{H^2} \simeq \begin{cases} 0.07 \times \left(\frac{10 T_{N_1}}{M_1}\right)^2 \left(\frac{M_1}{H}\right)^{3/2} \left(\frac{T_R}{T_{R_C}}\right) & \text{for } T_R < T_{R_C}, \\ 0.03 \times \left(\frac{10 T_{N_1}}{M_1}\right)^2 \left(\frac{M_1}{H}\right) & \text{for } T_R > T_{R_C}, \end{cases} \tag{22}$$

where we have used the relation given in Eq. (13). Therefore, we can safely neglect the above thermal effect, as long as M_1 is a bit larger than T_{N_1}. Next, there is another thermal effect which has been pointed out in Ref. [29]. If the temperature is lower than the effective mass for L and H_u, $T < m_{\mathrm{eff}} = \sqrt{\Sigma_\alpha |h_{1\alpha}|^2} |\tilde{N}_1|$, the evolution of the running gauge and/or Yukawa coupling constants $f(\mu)$ which couple to them are modified below the scale $\mu = m_{\mathrm{eff}}$. Thus, these running cou-

[8] One might wonder if the decay of the next-to-lightest SUSY particle into gravitino during or after the BBN would spoil the success of the BBN in the GMSB scenario. However, this problem is avoided for $m_{3/2} \lesssim 1$ GeV [23,24].

pling constants depend on $|\tilde{N}_1|$, and there appears an additional thermal potential for $\tilde{N}_1$:

$$\delta V(\tilde{N}_1) = aT^4 \log\left(\frac{|\tilde{N}_1|^2}{T^2}\right), \tag{23}$$

where a is a constant of order $\mathcal{O}(f^4)$. However, again, it turns out that the effective thermal mass for $\tilde{N}_1$ is less than the Hubble expansion rate

$$\frac{\delta M'^2_{\rm th}}{H^2} = \frac{aT^4}{H^2|\tilde{N}_1|^2} \simeq \begin{cases} 0.2\times a\left(\frac{M_{pl}}{|\tilde{N}_1|}\right)^2\left(\frac{M_1}{H}\right)\left(\frac{T_R}{T_{R_C}}\right)^2 & \text{for } T_R < T_{R_C}, \\ 0.05\times a\left(\frac{M_{pl}}{|\tilde{N}_1|}\right)^2 & \text{for } T_R > T_{R_C}, \end{cases} \tag{24}$$

and hence this thermal effect is also irrelevant to the present scenario.

III. INITIAL AMPLITUDE

In the previous section, we have assumed that the initial amplitude of the $\tilde{N}_1$'s oscillation is $|\tilde{N}_{1i}| \simeq M_{pl}$. This can be realized when the right-handed neutrino has only the mass term up to the Planck scale. In this section, we discuss another possibility, where the masses of the right-handed neutrinos are dynamically induced by a spontaneously broken gauge symmetry. The simplest candidate is $U(1)_{B-L}$, where B and L are baryon and lepton number, respectively. Let us denote the chiral superfields whose vacuum expectation values break the $U(1)_{B-L}$ by Φ and $\bar{\Phi}$. [We need two fields with opposite charges in order to cancel $U(1)_{B-L}$ gauge anomalies.] Because of the D term and F terms coming from the superpotential which gives the right-handed neutrino masses, the scalar potential of the right-handed sneutrino $\tilde{N}_1$ is lifted above the $U(1)_{B-L}$ breaking scale $\langle\Phi\rangle$ [5]. Therefore, the initial amplitude of the $\tilde{N}_1$'s oscillation at $H \simeq M_1$ is given by $|\tilde{N}_{1i}| \sim \langle\Phi\rangle$.

The breaking scale of the $U(1)_{B-L}$ gauge symmetry is model dependent. If it is broken at the Planck scale $\langle\Phi\rangle \simeq M_{pl}$, the discussion in the previous section does not change at all.[9] On the other hand, if $\langle\Phi\rangle$ is below the Planck scale, the initial amplitude of the $\tilde{N}_1$'s oscillation is reduced, and some parts of the discussion in the previous section are modified; those are the condition of the $\tilde{N}_1$ dominant universe [Eq. (14)] and the effective temperature of the cosmological gravitino problem [Eq. (17)].[10] [Notice that the amount of the generated lepton asymmetry given in Eq. (9) does not depend on the initial amplitude $|\tilde{N}_{1i}|$ as long as the $\tilde{N}_1$ dominant universe is realized.] Let us take $|\tilde{N}_{1i}| \sim \langle\Phi\rangle \sim 10^{17}$ GeV, for example. Due to the reduced initial amplitude of $\tilde{N}_1$, which means a smaller initial energy density, the condition for $\tilde{N}_1$ to dominate the universe is now given by

$$T_R \gg T_{N_1}\left(\frac{\sqrt{3}M_{pl}}{|\tilde{N}_{1i}|}\right)^2 \simeq 2\times 10^9 - 2\times 10^{10}\ \text{GeV}\times\left(\frac{|\tilde{N}_{1i}|}{10^{17}\ \text{GeV}}\right)^{-2}. \tag{25}$$

This condition is still easily satisfied by considering an inflation with relatively high scale. On the other hand, the effective temperature for the gravitino problem now becomes

$$T_{R\,\rm eff} = \begin{cases} T_{N_1}\left(\frac{\sqrt{3}M_{pl}}{|\tilde{N}_{1i}|}\right)^2, \\ T_{N_1}\left(\frac{\sqrt{3}M_{pl}}{|\tilde{N}_{1i}|}\right)^2\left(\frac{T_R}{T_{R_C}}\right), \end{cases} = \begin{cases} 2\times 10^9 - 2\times 10^{10}\ \text{GeV}\times\left(\frac{|\tilde{N}_{1i}|}{10^{17}\ \text{GeV}}\right)^{-2} & (\text{for } T_R < T_{R_C}), \\ 2\times 10^9 - 2\times 10^{10}\ \text{GeV}\times\left(\frac{|\tilde{N}_{1i}|}{10^{17}\ \text{GeV}}\right)^{-2}\left(\frac{T_R}{T_{R_C}}\right) & (\text{for } T_R > T_{R_C}). \end{cases} \tag{26}$$

[9] In this case, we need small couplings in order to explain the intermediate right-handed neutrino mass scale. For example, a superpotential $W = (1/2)y_i\Phi N_i N_i$ with $\langle\Phi\rangle \simeq M_{pl}$ and $y_3 \sim 10^{-4}$ gives the mass $M_3 \sim 10^{14}$ GeV to the heaviest right-handed neutrino. Such a small Yukawa coupling could well be a consequence of broken flavor symmetries.

[10] For the reduced initial amplitude, the thermal effect from the $aT^4\log(|\tilde{N}_1|^2)$ potential becomes larger than the case of $|\tilde{N}_{1i}| \simeq M_{pl}$. However, it is still irrelevant for $a \sim \mathcal{O}(f^4) \lesssim 10^{-2}$, as long as $|\tilde{N}_{1i}| \gtrsim 10^{17}$ GeV. [See Eq. (24).]

Thus, in this case, when the gravitino is unstable, its mass should be in a range of $m_{3/2} \gtrsim 1$ TeV to avoid the cosmological gravitino problem. This difficulty might be avoided when the gravitino is stable with mass $m_{3/2} \sim 10-100$ GeV [24].

IV. DISCUSSION AND CONCLUSIONS

We have investigated in this paper leptogenesis from the universe dominated by the right-handed sneutrino. We have found that this scenario is very successful in explaining the present baryon asymmetry. It is interesting that the amount of the generated lepton asymmetry is determined mainly by the decay temperature of the right-handed neutrino, independently of the reheating temperature T_R of the inflation. The desirable amount of the baryon asymmetry in the present universe is obtained when the decay temperature of the right-handed neutrino is $T_{N_1} \simeq 10^6 - 10^7$ GeV.

An attractive feature of this scenario is the entropy production by the decay of the coherent right-handed sneutrino, which itself produces the lepton asymmetry. The abundance of the thermally produced gravitinos is diluted by this entropy production, and the cosmological gravitino problem can be avoided in a wide range of the gravitino mass $m_{3/2} \simeq 10$ MeV–100 TeV. Actually, we have shown that the effective temperature $T_{R\,\mathrm{eff}}$, to which the constraint from the gravitino problem is applied, can be as low as $T_{R\,\mathrm{eff}} \simeq 2\times10^6 - 2\times10^7$ GeV, even with such high reheating temperatures as $T_R \gg 10^{10}$ GeV. The fact that such a high reheating temperature is allowed is very welcome from the viewpoint of building SUSY inflation models.

In particular, if the gravitino is stable, as in the gauge-mediated SUSY breaking models, the present energy density of the gravitino is determined by the decay temperature of the right-handed neutrino $T_{N_1} \simeq 10^6 - 10^7$ GeV (if we assume the initial amplitude of the coherent right-handed sneutrino is $|\tilde{N}_{1i}| \simeq M_{pl}$). Thus, the gravitino mass can be predicted from the observed energy density of the dark matter as $m_{3/2} \simeq 10$ MeV-1 GeV, for a wide range of the reheating temperature 10^6 GeV$< T_R < 7\times10^{11}(M_1/10^6\ \mathrm{GeV})^{1/2}$ GeV, assuming that the dark matter in our universe is dominantly composed of the gravitino.

Finally, we comment on the isocurvature density perturbation coming from the fluctuation of the initial amplitude of the right-handed sneutrino, $|\tilde{N}_{1i}|$.[11] The baryonic isocurvature perturbation from $\tilde{N}_1$ is given by

$$\delta_{\mathrm{iso}} = \frac{\delta n_B^{\mathrm{iso}}}{n_B}\frac{\Omega_B}{\Omega_t} \simeq \frac{H_{\mathrm{inf}}}{\pi|\tilde{N}_{1i}|}\frac{\Omega_B}{\Omega_t} \simeq 1\times10^{-7}\left(\frac{M_{pl}}{|\tilde{N}_{1i}|}\right)\left(\frac{H_{\mathrm{inf}}}{10^{13}\ \mathrm{GeV}}\right)\left(\frac{\Omega_B}{0.1\times\Omega_t}\right), \tag{27}$$

where Ω_B and Ω_t is the density parameters of baryons and total matter, respectively. This isocurvature fluctuation might be detected in future experiments.

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APPENDIX

In the body of this paper, we have discussed the leptogenesis scenario from the universe dominated by $\tilde{N}_1$. Here, we briefly comment on the case where the $\tilde{N}_1$ does not dominate the universe. In this case, the resultant lepton asymmetry depends on the reheating temperature T_R and the initial amplitude of the oscillation $|\tilde{N}_{1i}|$, and it is given in the following form [5]:

$$\frac{n_L}{s} = \frac{1}{4}\epsilon_1\left(\frac{T_R}{M_1}\right)\left(\frac{|\tilde{N}_{1i}|}{\sqrt{3}M_{pl}}\right)^2 \cong 0.8\times10^{-10}\left(\frac{T_R}{10^7\mathrm{GeV}}\right)\left(\frac{|\tilde{N}_{1i}|}{M_{pl}}\right)^2\left(\frac{m_{\nu_3}}{0.05\ \ \mathrm{eV}}\right)\delta_{\mathrm{eff}}. \tag{A1}$$

Thus, it is possible to produce the desired amount of baryon asymmetry, avoiding the cosmological gravitino problem, although it depends crucially on the reheating temperature.

[1] Super-Kamiokande Collaboration, Y. Fukuda *et al.*, Phys. Lett. B **433**, 9 (1998); **436**, 33 (1998); Phys. Rev. Lett. **81**, 1562 (1998); see also, Super-Kamiokande Collaboration, C. McGrew, talk presented at The 2nd International Workshop on Neutrino Oscillations and their Origin ("NOON2000"), Tokyo, Japan, 2000.

[2] T. Yanagida, Prog. Theor. Phys. **64**, 1103 (1980); in *Proceedings of the Workshop on the Unified Theory and the Baryon*

Number in the Universe, Tsukuba, Japan, 1979, edited by O. Sawada and A. Sugamoto, (KEK Report No.-79-18, Tsukuba, 1979), p 95; M. Gell-Mann, P. Ramond, and R. Slansky, in *Supergravity*, edited by D. Z. Freedman and P. van Nieuwenhuizen (North-Holland, Amsterdam, 1979).

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

[4] V. A. Kuzmin, V. A. Rubakov, and M. E. Shaposhnikov, Phys. Lett. **155B**, 36 (1985).

[5] H. Murayama and T. Yanagida, Phys. Lett. B **322**, 349 (1994).

[6] T. Moroi and H. Murayama, J. High Energy Phys. **07**, 009 (2000).

[7] T. Asaka, M. Fujii, K. Hamaguchi, and T. Yanagida, Phys. Rev. D **62**, 123514 (2000); M. Fujii, K. Hamaguchi, and T. Yanagida, *ibid.* **63**, 123513 (2001).

[8] I. Affleck and M. Dine, Nucl. Phys. **B249**, 361 (1985).

[9] M. Y. Khlopov and A. D. Linde, Phys. Lett. **138B**, 265 (1984); J. Ellis, J. E. Kim, and D. V. Nanopoulos, *ibid.* **145B**, 181 (1984); M. Kawasaki and T. Moroi, Prog. Theor. Phys. **93**, 879 (1995); see also recent analyses, E. Holtmann, M. Kawasaki, K. Kohri, and T. Moroi, Phys. Rev. D **60**, 023506 (1999); M. Kawasaki, K. Kohri, and T. Moroi, *ibid.* **63**, 103502 (2001).

[10] M. H. Reno and D. Seckel, Phys. Rev. D **37**, 3441 (1988); see also a recent analysis, K. Kohri, *ibid.* **64**, 043515 (2001).

[11] H. Pagels and J. R. Primack, Phys. Rev. Lett. **48**, 223 (1982).

[12] T. Moroi, H. Murayama, and M. Yamaguchi, Phys. Lett. B **303**, 289 (1993); A. de Gouvêa, T. Moroi, and H. Murayama, Phys. Rev. D **56**, 1281 (1997).

[13] M. Dine, A. E. Nelson, and Y. Shirman, Phys. Rev. D **51**, 1362 (1995); M. Dine, A. E. Nelson, Y. Nir, and Y. Shirman, *ibid.* **53**, 2658 (1996); for a review, see, for example, G. F. Giudice and R. Rattazzi, Phys. Rep. **322**, 419 (1999).

[14] Y. Nomura and T. Yanagida, Phys. Lett. B **487**, 140 (2000).

[15] M. Dine, L. Randall, and S. Thomas, Phys. Rev. Lett. **75**, 398 (1995); Nucl. Phys. **B458**, 291 (1996).

[16] M. K. Gaillard, H. Murayama, and K. A. Olive, Phys. Lett. B **355**, 71 (1995).

[17] L. Covi, E. Roulet, and F. Vissani, Phys. Lett. B **384**, 169 (1996); M. Flanz, E. A. Paschos, and U. Sarkar, *ibid.* **345**, 248 (1995); W. Buchmüller and M. Plümacher, *ibid.* **431**, 354 (1998).

[18] W. Buchmuller and T. Yanagida, Phys. Lett. B **445**, 399 (1999).

[19] S. Y. Khlebnikov and M. E. Shaposhnikov, Nucl. Phys. **B308**, 885 (1988); J. A. Harvey and M. S. Turner, Phys. Rev. D **42**, 3344 (1990).

[20] See, for example, E. Kolb and M. Turner, *The Early Universe* (Addison-Wesley, New York, 1990).

[21] H. Murayama, H. Suzuki, T. Yanagida, and J. Yokoyama, Phys. Rev. Lett. **70**, 1912 (1993); Phys. Rev. D **50**, 2356 (1994).

[22] See, for example, T. Asaka, K. Hamaguchi, M. Kawasaki, and T. Yanagida, Phys. Rev. D **61**, 083512 (2000), and references therein; M. Kawasaki, M. Yamaguchi, and T. Yanagida, Phys. Rev. Lett. **85**, 3572 (2000).

[23] Moroi, Murayama, and Yamaguchi [12].

[24] M. Bolz, W. Buchmüller, and M. Plümacher, Phys. Lett. B **443**, 209 (1998); T. Gherghetta, G. F. Giudice, and A. Riotto, *ibid.* **446**, 28 (1999); T. Asaka, K. Hamaguchi, and K. Suzuki, *ibid.* **490**, 136 (2000).

[25] Particle Data Group, D. E. Groom *et al.*, Eur. Phys. J. C **15**, 1 (2000).

[26] Fujii, Hamaguchi, and Yanagida [7].

[27] J. Sato and T. Yanagida, Nucl. Phys. B (Proc. Suppl.) **77**, 293 (1999); P. Ramond, *ibid.* **77**, 3 (1999).

[28] R. Allahverdi, B. A. Campbell, and J. Ellis, Nucl. Phys. **B579**, 355 (2000).

[29] A. Anisimov and M. Dine, hep-ph/0008058.