Abstract

We study the cosmic microwave background (CMB) anisotropy in the scenario where the baryon asymmetry of the universe is generated from a condensation of a scalar field. In such a scenario, the scalar condensation may acquire fluctuation during the inflation which becomes a new source of the cosmic density perturbations. In particular, the primordial fluctuation of the scalar condensation may induce correlated mixture of the adiabatic and isocurvature fluctuations. If the scalar condensation decays before it completely dominates the universe, the CMB angular power spectrum may significantly deviate from the conventional adiabatic result. Such a deviation may be observed in the on-going MAP experiment.
1 Introduction

In modern particle cosmology, one of the most important issues is to understand the origin of the baryon asymmetry of the universe. In particular, assuming inflation [1] as a solution to the horizon, flatness, and other cosmological problems and as a seed for density fluctuations for later structure formation, we are obliged to adopt scenarios where the baryon asymmetry is generated after the reheating.

One of the main branches of baryogenesis is to use a primordial condensation of a scalar field as a source of the baryon-number asymmetry of the universe. In such a class of scenarios, there exists a primordial condensation of a scalar field and it decays at a later stage of the evolution of the universe. Then, baryon asymmetry is generated by the decay of the scalar field. If such a scalar field exists, one of the Sakharov’s three conditions for baryogenesis, i.e., the out-of-equilibrium condition, is easily satisfied by assuming non-vanishing initial amplitude of the scalar field. In particular, in cosmological scenarios based on supersymmetric models, there exist various possible candidates of such a scalar field and many efforts have been made to study such scenarios.

Among various possibilities, probably one of the most well-motivated candidates of such scalar fields is the superpartner of the right-handed neutrino, the right-handed sneutrino. Assuming the seesaw mechanism [2] as a origin of very tiny neutrino masses which are suggested from the neutrino-oscillation experiments [3] the right-handed sneutrinos inevitably exist in the supersymmetric models. Such a right-handed sneutrino can be a source of the baryon-number asymmetry [4, 5]; if one of the right-handed neutrinos has non-vanishing amplitude in the early universe, its decay may generate lepton-number asymmetry which is converted to the baryon-number asymmetry via the sphaleron interaction [6]. (In fact, this is a supersymmetric extension of the leptogenesis scenario proposed by Fukugita and Yanagida [7].)

In addition, Affleck-Dine mechanism for baryogenesis [8] is another possibility. In the Affleck-Dine scenario, using the baryon-number violating operator as a source, non-vanishing baryon number is generated while a scalar partner of quarks, called Affleck-Dine field, is oscillating.

Since the baryogenesis by a scalar condensation is attractive and well-motivated in particular in supersymmetric models, it is important to consider how the scenario can be experimentally tested. As we will discuss below, one possible effect of such a scalar field is on the cosmological density perturbations; if a condensation of the scalar field in the early universe is the source of the present baryon asymmetry, some signal may be imprinted in the cosmic microwave background (CMB) since the significant fraction of the CMB radiation we observe today may also originate in the scalar-field condensation. In particular, the amplitude of the scalar field may acquire sizable fluctuation during the inflation and such a fluctuation affects the CMB anisotropy.

Thus, in this paper, we study the cosmic density perturbations in the scenario of the

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#1 Even for Dirac neutrinos, an analogue of the seesaw mechanism is possible [4, 5] and the leptogenesis is possible [6] via the decay of a scalar condensate [7].
baryogenesis by a scalar field. Although our conclusions are quite general to a large class of scenarios, we mostly concentrate on the scenario of the sneutrino-induced leptogenesis [8, 9] to make our discussion clearer. In particular, we study effects of the primordial fluctuation of the scalar field on the CMB anisotropy. If the primordial fluctuation exists in the amplitude of the scalar field, correlated mixture of the adiabatic and isocurvature fluctuations may be generated which may significantly affect the CMB angular power spectrum. With the precise measurement of the CMB angular power spectrum by the MAP experiment, some signal of the scalar-field-induced leptogenesis may be observed.

2 The Leptogenesis Scenario

Let us start by introducing the scenario we consider. For our argument, there are two scalar fields which play important roles; one is the inflaton field $\chi$ and the other is the (lightest) right-handed sneutrino. In the following, we denote the lightest right-handed scalar neutrino as $\tilde{N}$. We assume the inflation so that the universe starts with a de Sitter epoch when the universe is dominated by the potential energy of the inflaton field $\chi$. During the inflation, the right-handed sneutrino is assumed to have a non-vanishing amplitude $\tilde{N}_{\text{init}}$. We treat $\tilde{N}_{\text{init}}$ as a free parameter in the following discussion. After the inflation, the inflaton field starts to oscillate and then decays. The right-handed neutrino also starts to oscillate when the expansion rate of the universe becomes comparable to the Majorana mass $M_N$. (We assume that the expansion rate during the inflation is larger than $M_N$ so that $\tilde{N}$ starts to oscillate after the inflation.) Then, the right-handed sneutrino decays when $H \sim \Gamma_N$, where $\Gamma_N$ is the decay rate of $\tilde{N}$.

One important point in this scenario is that the CMB radiation we observe today has two origins; the inflaton field and the right-handed sneutrino. This is because the decays of the inflaton and the right-handed sneutrino both convert the energy densities stored in the scalar condensation into that of radiation. For our study, it is convenient to distinguish the radiation originating in the inflaton field from that originating in the right-handed sneutrino. We denote them as $\gamma_\chi$ and $\gamma_{\tilde{N}}$, respectively.\(^\#2\) Energy fraction of the radiation from $\tilde{N}$ depends on the initial amplitude $\tilde{N}_{\text{init}}$ as well as the decay rate of the inflaton $\chi$ (denoted as $\Gamma_\chi$), decay rate of $\tilde{N}$ and the mass of the right-handed (s)neutrino.

If $\Gamma_\chi \lesssim \Gamma_N$, $\tilde{N}$ decays when the universe is dominated by the inflaton condensation. In fact, we cannot neglect the momentum transfer between these photons and they cannot be defined separately. In particular, velocity perturbations of these photons should be the same since the mean free path of the photon is much shorter than the horizon scale. If we only discuss the behavior of $\delta_{\gamma_X}$ for the super-horizon modes, however, the following discussions are valid since the velocity perturbations are suppressed by the factor $k\tau$ relative to the density perturbations. In a rigorous sense, $\gamma_\chi$ and $\gamma_{\tilde{N}}$ should be understood as representatives of the components produced from the decay products of the inflaton field and $\tilde{N}$, respectively.

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particular, when $\Gamma_N \lesssim H \lesssim M_N$, energy densities of $\tilde{N}$ and $\chi$ are related as\(^\#3\)

$$\frac{\rho_{\tilde{N}}}{\rho_\chi} \sim \frac{\tilde{N}_{\text{init}}^2}{M_*^2} : \Gamma_N \lesssim H \lesssim M_N. \quad (2.1)$$

(Here and hereafter, $\rho_\chi$ denotes the energy density of the component $X$.) Using this relation, we can evaluate the energy density of the right-handed sneutrino at the time of its decay; when $H \sim \Gamma_N$, $\rho_\chi \sim \Gamma_N^2 M_*^2$ and hence $\rho_{\tilde{N}} \sim \Gamma_N^2 \tilde{N}_{\text{init}}^2$. After the decay, energy density stored in the sneutrino condensation is converted to that of radiation, so the universe is filled with the radiation (denoted as $\gamma_N$) as well as the inflaton oscillation. Using the relations $\rho_\chi \propto a^{-3}$ and $\rho_{\tilde{N}} \propto a^{-4}$, we obtain the energy densities of each component at $H \sim \Gamma_\chi$ as $\rho_\chi \sim \Gamma_\chi^2 M_*^2$ and $\rho_{\tilde{N}} \sim \rho_\chi \Gamma_\chi^2 M_*^2 (\tilde{N}_{\text{init}}/M_*^2)^2$. Thus, after the decay of the inflaton field, the energy fraction of the radiation from the right-handed sneutrino is estimated as

$$f_{\gamma_{\tilde{N}}} \sim \frac{\left(\Gamma_\chi/\Gamma_N\right)^{2/3}(\tilde{N}_{\text{init}}/M_*^2)^2}{1 + \left(\Gamma_\chi/\Gamma_N\right)^{2/3}(\tilde{N}_{\text{init}}/M_*^2)^2} : \Gamma_\gamma \lesssim \Gamma_N. \quad (2.2)$$

Here, $f_{\gamma_X}$ is the energy fraction of $\gamma_X$ after the decays of the inflaton and $\tilde{N}$, and $f_{\gamma_X} + f_{\gamma_{\tilde{N}}} = 1$. (Here and hereafter, $\gamma_X = \gamma_{\tilde{N}}$ or $\gamma_X$.)

If $\Gamma_\chi \gtrsim \Gamma_N$, on the contrary, $\tilde{N}$ decays after the inflaton decay. In this case, it is convenient to define the following quantity:

$$\tilde{N}_{\text{eq}} \sim \begin{cases} \left(\Gamma_N/M_*^2\right)^{1/4} : M_N < \Gamma_\chi \\ \left(\Gamma_N/\Gamma_\chi\right)^{1/4} : M_N > \Gamma_\chi \end{cases}. \quad (2.3)$$

If $\tilde{N}_{\text{init}} \sim \tilde{N}_{\text{eq}}$, $\rho_\chi \sim \rho_{\tilde{N}}$ is realized when $H \sim \Gamma_\tilde{N}$. Thus, when $\tilde{N}_{\text{init}} \lesssim \tilde{N}_{\text{eq}}$, $\tilde{N}$ decays in the $\gamma_X$-dominated universe and hence

$$f_{\gamma_{\tilde{N}}} \sim \frac{(\tilde{N}_{\text{init}}/\tilde{N}_{\text{eq}})^2}{1 + (\tilde{N}_{\text{init}}/\tilde{N}_{\text{eq}})^2} : \Gamma_N \lesssim \Gamma_\chi, \quad \tilde{N}_{\text{init}} \lesssim \tilde{N}_{\text{eq}}. \quad (2.4)$$

On the contrary, if $\tilde{N}_{\text{init}} \gtrsim \tilde{N}_{\text{eq}}$, the right-handed sneutrino decays after it dominates the universe and we obtain

$$f_{\gamma_{\tilde{N}}} \sim \frac{(\tilde{N}_{\text{init}}/\tilde{N}_{\text{eq}})^{8/3}}{1 + (\tilde{N}_{\text{init}}/\tilde{N}_{\text{eq}})^{8/3}} : \Gamma_N \lesssim \Gamma_\chi, \quad \tilde{N}_{\text{eq}} \lesssim \tilde{N}_{\text{init}}. \quad (2.5)$$

\(^\#3\)When the initial amplitude of $\tilde{N}$ is larger than $M_*$, inflation is caused by the energy density of $\tilde{N}$. If this happens, all the components generated from $\chi$ is washed out and $f_{\gamma_{\tilde{N}}} \rightarrow 1$. If the e-folding number of the secondary inflation caused by $\tilde{N}$ is small enough, the following discussions are unchanged since the density fluctuation for the scale corresponding to the CMB anisotropy we will discuss is generated from the primary inflation caused by $\chi$. If the initial amplitude is much larger than $M_*$, only the energy density from the decay of the right-handed neutrino is relevant and the perturbation becomes purely adiabatic.\[\]
When $\tilde{N}$ decays, lepton-number asymmetry is also generated. Such a lepton-number asymmetry is converted to the baryon-number asymmetry due to the sphaleron process. The resultant baryon-number asymmetry depends on whether $\tilde{N}$ decays before or after the $\tilde{N}$-dominated universe is realized. If $\tilde{N}$ decays later than the inflaton decay, the baryon-to-entropy ratio is given by

$$n_B/s \simeq 0.24 \times 10^{-10} f_{\gamma_N} \delta_{\text{eff}} \left( T_N/10^6 \text{GeV} \right) \left( m_{\nu_3}/0.05 \text{eV} \right) : \Gamma_N < \Gamma_\chi,$$

where $T_N$ is the temperature at the epoch of the decay of $\tilde{N}$, $m_{\nu_3}$ the mass of the heaviest (left-handed) neutrino mass. (So, if $\tilde{N}$ decays after dominating the universe, $T_N$ becomes the reheat temperature due to the decay of $\tilde{N}$.) In addition, in the basis where the Majorana mass matrix for the right-handed neutrinos $\hat{M}$ is real and diagonalized, the effective CP violating phase is given by

$$\delta_{\text{eff}} = \frac{\langle H_u \rangle^2 \text{Im}[\hat{h}^\dagger \hat{M}^{-1} \hat{h}^* \hat{h}^T]_{11}}{[h h^\dagger]_{11}},$$

where $\hat{h}_{i\alpha}$ is the neutrino Yukawa matrix with $i$ and $\alpha$ being the generation indices for the right-handed and left-handed neutrinos, respectively. Notice that, with a maximum CP violation, $\delta_{\text{eff}} \sim 1$. On the contrary, if $\Gamma_\chi < \Gamma_N$, $\tilde{N}$ decays before the decay of the inflaton field. In this case, radiation-dominated universe is realized much later than the epoch of the sneutrino decay. In this case, the baryon-to-entropy ratio is given by

$$n_B/s \simeq 0.24 \times 10^{-10} f_{\gamma_{\tilde{N}}} \delta_{\text{eff}} \left( T_R/10^6 \text{GeV} \right) \left( m_{\nu_3}/0.05 \text{eV} \right) \left( \tilde{N}_{\text{init}}/M_* \right)^2 : \Gamma_\chi < \Gamma_N,$$

where $T_R$ is the reheating temperature due to the decay of the inflaton field. As one can see, large enough baryon asymmetry ($n_B/s \sim 8 \times 10^{-11}$) can be generated in this scenario with relatively low reheating temperature. In particular, even if we impose the upper bound on the reheating temperature to avoid the overproduction of the gravitinos, it is possible to generate large enough baryon asymmetry. For $\Gamma_N < \Gamma_\chi$, requiring $T_N \lesssim 10^9-10^{10}$ GeV in order not to overproduce gravitinos [3], $n_B/s \sim 10^{-10}$ can be realized when $f_{\gamma_{\tilde{N}}} \gtrsim 10^{-3} - 10^{-2}$. For $\Gamma_\chi < \Gamma_N$, much smaller $f_{\gamma_{\tilde{N}}}$ is possible to generate enough baryon asymmetry even if we require $T_R \lesssim 10^9-10^{10}$ GeV to avoid the gravitino problem. In addition, it is notable that $n_B/s \sim 10^{-10}$ can be realized even with $f_{\gamma_{\tilde{N}}} \ll 1$ if we push up $T_N$ or $T_R$.

### 3 Density Perturbations

As we have seen, in the scenario we consider, baryon asymmetry as well as some fraction of the CMB radiation are generated from the decay product of $\tilde{N}$. Thus, if $\tilde{N}$ has a primordial fluctuation, such a fluctuation also becomes the source of the cosmic density
perturbations. Most importantly, such a fluctuation affects the CMB anisotropy which is now being measured very precisely with various experiments [16, 17, 18, 19, 20].

Fluctuation of \( \tilde{N} \) is primarily induced during the inflation; assuming that the (effective) mass of the right-handed sneutrino during the inflation is smaller than the expansion rate during the inflation \( H_{\text{inf}} \), the primordial fluctuation of \( \tilde{N} \) is estimated as

\[
\delta \tilde{N}_{\text{init}} = \frac{H_{\text{inf}}}{2\pi}.
\] (3.1)

If \( \delta \tilde{N}_{\text{init}} \) is non-vanishing, this becomes a source of the entropy between components generated from the decay products of \( \chi \) and \( \tilde{N} \). It is convenient to define the parameter

\[
S^{(\delta \tilde{N})}_{\tilde{N}\chi} \equiv \delta^{(\delta \tilde{N})}_{\tilde{N}} - \delta^{(\delta \tilde{N})}_{\chi},
\] (3.2)

where \( \delta_X \equiv \delta \rho_X / \rho_X \) with \( \rho_X \) being the energy density of the component \( X \) and \( \delta \rho_X \) its perturbation in the Newtonian gauge. Here, the right-hand side is evaluated after \( \tilde{N} \) and \( \chi \) both start to oscillate, and the superscript “\((\delta \tilde{N})\)” is for variables generated from the primordial fluctuation of \( \tilde{N} \). Solving the equations of motions for the scalar fields, we obtain [21, 22]

\[
S^{(\delta \tilde{N})}_{\tilde{N}\chi} = \frac{2\delta \tilde{N}_{\text{init}}\tilde{N}_{\text{init}}}{N_{\text{init}}},
\] (3.3)

where we assumed that the potential of the scalar fields are dominated by the quadratic terms. After the decays of \( \tilde{N} \) and \( \chi \), \( S^{(\delta \tilde{N})}_{\tilde{N}\chi} \) becomes entropy between components generated from the decay products of \( \tilde{N} \) and \( \chi \) for the perturbations with wavelength much longer than the horizon scale.

With the non-vanishing value of the primordial fluctuation of \( \tilde{N} \) given in Eq. (3.1), it is important to note that there are two independent sources of the perturbations in our case; fluctuation of the inflaton field and that of the right-handed sneutrino. In the framework of the linear perturbation theory, we can discuss effects of these perturbations separately. In addition, in discussing the CMB angular power spectrum \( C_l \), effects of the two perturbations can be treated separately since we assume no correlation between these two fields. Thus, the total angular power spectrum is given in the form

\[
C_l = C^{(\delta \chi)}_l + C^{(\delta \tilde{N})}_l,
\] (3.4)

where \( C^{(\delta \chi)}_l \) and \( C^{(\delta \tilde{N})}_l \) are the contributions from the primordial fluctuations in the inflaton and \( \tilde{N} \), respectively. The inflaton contribution \( C^{(\delta \chi)}_l \) is known to become the adiabatic result (with relevant spectral power index). Thus, let us consider the second term \( C^{(\delta \tilde{N})}_l \).

In discussing the cosmic density perturbations induced from \( \delta \tilde{N}_{\text{init}} \), we treat the radiations originated from these fields separately; \( \gamma_{\chi} \) for the radiation from the inflaton and
which is that from the right-handed sneutrino. Density perturbations for these components are defined separately: \( \delta_{\gamma_X} \equiv \delta \rho_{\gamma_X} / \rho_{\gamma_X} \) and \( V_{\gamma_X} \) are density and velocity fluctuations of \( \gamma_X \) with \( X \) being \( \chi \) and \( \tilde{N} \). (Hereafter, we use the Newtonian gauge. We follow the notation and convention of [23].)

In the following, we follow the evolution of the density perturbations of various components solving the relevant Einstein and Boltzmann equations. In particular, we are interested in evolutions of the density perturbations when \( H \gtrsim \Gamma_N \) and properties of the density perturbations just after the decay of \( \tilde{N} \). In this case, we consider the evolution of the perturbations in the universe with a very high temperature where various charged particles are thermally produced. Then, the radiation component becomes locally isotropic and hence the anisotropic stress perturbation can be neglected.

Denoting the perturbed line element as

\[
\begin{aligned}
ds^2 &= a^2 \left[ -(1 + 2\Psi) d\tau^2 + (1 + 2\Phi) \delta_{ij} dx^i dx^j \right],
\end{aligned}
\]

(3.5)

with \( \tau \) being the conformal time and \( a \) being the scale factor, the Poisson equation for the metric perturbation is given by

\[
\begin{aligned}
k^2 \Phi &= 4\pi G a^2 \rho_T \left[ \delta_T + 3 \frac{a'}{a} (1 + w_T) V_T / k \right],
\end{aligned}
\]

(3.6)

where \( k \) is the comoving momentum, \( G \) the Newton constant, and the “prime” denotes the derivative with respect to the conformal time \( \tau \). In addition, the subscript “\( T \)” is for the total matter and \( w_T \) denotes the equation-of-state parameter of the total matter. On the other hand, neglecting the contribution from the anisotropic stress tensor, evolution of the density and velocity perturbations in the Fourier space are given by [23]

\[
\begin{aligned}
\delta'_{\gamma_X} &= -\frac{4}{3} k V_{\gamma_X} - 4\Phi', \\
V'_{\gamma_X} &= k \left( \frac{1}{4} \delta_{\gamma_X} + \Psi \right).
\end{aligned}
\]

(3.7)

In addition, in the situation we consider, the relation \( \Phi = -\Psi \) holds since the anisotropic stress tensor is small enough. In the following, we use this relation to eliminate \( \Phi \).

Using Eqs. (3.6) and (3.7), we can discuss the evolution of the cosmic density perturbations induced from the primordial fluctuation of the amplitude of the right-handed sneutrino condensation. In this case, we neglect the adiabatic fluctuations generated by the inflaton fluctuation. In addition, since we are interested in the perturbations whose wavelength is much longer than the horizon scale (at the epoch of the baryogenesis and reheating), we can expand the solution as a function of \( k\tau \).

Evolution of \( \delta_{\gamma_X} \) is quite simple. Just after the inflation, the right-handed sneutrino is the sub-dominant component and hence \( \Psi^{(\delta \tilde{N})} \) vanishes as \( \tau \to 0 \). In addition, since there is no primordial perturbation in the inflaton sector for this mode, \( \delta_{\gamma_X}^{(\delta \tilde{N})} \to 0 \) as \( \tau \to 0 \). Thus, we obtain

\[
\delta_{\gamma_X}^{(\delta \tilde{N})} = 4\Psi^{(\delta \tilde{N})} + O(k^2 \tau^2).
\]

(3.8)
This relation holds at any moment of the evolution of the universe.

Using the above relation, the entropy between the radiation and the baryon components can be evaluated. We calculate the entropy in the radiation-dominated universe after the decay of both $\chi$ and $\tilde{N}$. In the radiation-dominated universe [23]

$$\Delta_T = O(k^2 \tau^2), \quad \delta_T = -2 \Psi_{\text{RD}}, \quad V_T = \frac{1}{2} \Psi_{\text{RD}} k \tau. \quad (3.9)$$

(Here and hereafter, we only denote the leading contribution to the perturbations.) Then, using the relation

$$\Delta_T = f_{\gamma \chi} \Delta_{\gamma \chi} + f_{\gamma \tilde{N}} \Delta_{\gamma \tilde{N}}, \quad (3.10)$$

with $\Delta_{\gamma \chi} = \delta_{\gamma \chi} + 4(a'/a)V_T/k$, we obtain

$$\Delta_{\gamma \tilde{N}} = \frac{f_{\gamma \chi}}{f_{\gamma \tilde{N}}} \Delta_{\gamma \chi} = -6 \frac{f_{\gamma \chi}}{f_{\gamma \tilde{N}}} \Psi_{\text{RD}}^{(\delta \tilde{N})}. \quad (3.11)$$

Using the fact that the entropy between any component produced from $\tilde{N}$ and that from the inflaton field is conserved, we can relate $\Psi_{\text{RD}}^{(\delta \tilde{N})}$ with $S_{\tilde{N} \chi}^{(\delta \tilde{N})}$; with relation $S_{\tilde{N} \chi}^{(\delta \tilde{N})} = \frac{3}{4} (\Delta_{\tilde{N}} - \Delta_{\tilde{N}})$, we obtain

$$\Psi_{\text{RD}}^{(\delta \tilde{N})} = - \frac{2}{9} f_{\gamma \tilde{N}} S_{\tilde{N} \chi}^{(\delta \tilde{N})} = - \frac{4}{9} f_{\gamma \chi} \frac{\delta \tilde{N}_{\text{init}}}{N_{\text{init}}}. \quad (3.12)$$

Thus, if $\tilde{N}$ decays after it dominates the universe, $f_{\gamma \tilde{N}} \approx 1$ and hence the metric perturbation becomes comparable to the primordial entropy perturbation. On the other hand, if $f_{\gamma \tilde{N}} \ll 1$, the metric perturbation becomes negligibly small.

Since the baryon asymmetry is generated from $\tilde{N}$, the density fluctuation in the baryonic component is given by

$$\Delta_{\tilde{N}} = \frac{3}{4} \Delta_{\gamma \tilde{N}} = - \frac{9}{2} f_{\gamma \chi} \Psi_{\text{RD}}^{(\delta \tilde{N})}. \quad (3.13)$$

Thus, the entropy between the baryon and the radiation is given by

$$S_{b \gamma}^{(\delta \tilde{N})} = \Delta_{\tilde{N}} - \frac{3}{4} \Delta_{\gamma \tilde{N}} = - \frac{9}{2} f_{\gamma \chi} \Psi_{\text{RD}}^{(\delta \tilde{N})} = - \frac{9(1 - f_{\gamma \tilde{N}})}{2 f_{\gamma \tilde{N}}} \Psi_{\text{RD}}^{(\delta \tilde{N})}. \quad (3.14)$$

Thus, the correlated mixture of the adiabatic and isocurvature fluctuations is generated. In addition, if the right-handed sneutrino once dominates the universe, $f_{\gamma \tilde{N}} \rightarrow 1$ and hence the perturbation becomes adiabatic [21, 23, 24, 25]. (Thus, in this case, $\tilde{N}$ may play the role of the so-called “curvaton” field.) On the contrary, if the inflaton field and its decay products always dominate the universe, $f_{\gamma \tilde{N}} \rightarrow 0$ and the perturbation becomes isocurvature. Thus, the interesting spectrum would be obtained when the inflaton and the right-handed sneutrino both produce significant amount of the resultant radiation.
4 CMB Angular Power Spectrum

Now we discuss the CMB angular power spectrum induced from the primordial fluctuation of $N$. For this purpose, we parameterize the entropy perturbation induced from $\delta N_{\text{init}}$ as

$$S_{\ell \gamma}^{(\delta N)} = \kappa_b \Psi_{\text{RD}}^{(\delta N)}.$$  \hspace{1cm} (4.1)

Then, using Eqs. (2.2), (2.4), (2.5), and (3.14), the $\kappa_b$ parameter is estimated as

$$\kappa_b \sim \frac{9}{2} \times \begin{cases} (\Gamma_\chi / \Gamma_N)^{-2/3} (\bar{N}_{\text{init}} / M_*)^{-2} & : \Gamma_\chi \ll \Gamma_N \end{cases}$$

$$\begin{cases} (\bar{N}_{\text{init}} / \bar{N}_{\text{eq}})^{-2} & : \Gamma_N \ll \Gamma_\chi, \bar{N}_{\text{init}} \ll \bar{N}_{\text{eq}} \end{cases}$$

$$\begin{cases} (\bar{N}_{\text{init}} / \bar{N}_{\text{eq}})^{-8/3} & : \Gamma_N \ll \Gamma_\chi, \bar{N}_{\text{eq}} \ll \bar{N}_{\text{init}} \end{cases}$$ \hspace{1cm} (4.2)

Thus, in our case, the $\kappa_b$ parameter varies from $-\infty$ to 0. If $\bar{N}_{\text{init}} \ll \bar{N}_{\text{eq}}$, the $\kappa_b$ parameter goes to $-\infty$ and hence the density perturbation becomes (baryonic) isocurvature. On the contrary, when $\bar{N}_{\text{init}} \gg \bar{N}_{\text{eq}}$, $\kappa_b \sim 0$ and the adiabatic density perturbation is obtained from the primordial fluctuation of $\bar{N}$. The most interesting case is $\bar{N}_{\text{init}} \sim \bar{N}_{\text{eq}}$, where $\kappa_b \sim O(0.1-1)$. In this case, correlation between adiabatic and isocurvature perturbations becomes the most effective. For the case where $\Gamma_N \ll \Gamma_\chi$, the requirement of enough baryon asymmetry with the gravitino constraint give $f_{\tilde{S}} \gtrsim 10^{-3} - 10^{-2}$, and hence $\kappa_b \gtrsim -1000$. If $\Gamma_\chi < \Gamma_N$, smaller $\kappa_b$ is possible. (See Eqs. (2.4) and (2.8).)

In Fig. 4, we plot the CMB angular power spectrum induced from the primordial fluctuation of the right-handed sneutrino $C_i^{(\delta N)}$. (Here and hereafter, we assume there is no entropy perturbation in the CDM sector.) Notice that the lines with $\kappa_b = 0$ and $\kappa_b = -\infty$ coincide with the results with purely adiabatic and isocurvature density perturbations, respectively. For a general case, however, the angular power spectrum has a unique structure. In particular, as one can see, the acoustic peaks are suppressed relative to the Sachs-Wolfe (SW) tail as $|\kappa_b|$ increases.

As we discussed, the total angular power spectrum is given by the sum of the inflaton and the $\bar{N}$ contributions, as shown in Eq. (3.4). Here, $C_i^{(\delta \chi)}$ is the inflaton contribution which is parameterized by the primordial metric perturbation generated from the inflaton perturbation $\Psi^{(\delta \chi)}$, which is given by [20]

$$\Psi_{\text{RD}}^{(\delta \chi)} = \frac{4}{9} \left[ \frac{H_{\text{inf}}}{2\pi} \frac{3H_{\text{inf}}^2}{V_{\text{inf}}^k} \right]_{k=aH_{\text{inf}}} \hspace{1cm} (4.3)$$

where $V_{\text{inf}}$ is the inflaton potential and $V_{\text{inf}}' \equiv (\partial V_{\text{inf}} / \partial \chi)$, and the superscript $(\delta \chi)$ means that the corresponding variable is generated from the inflaton fluctuation. On the contrary, $C_i^{(\delta N)}$ is from the primordial fluctuation of $\bar{N}$, and is parameterized by $S_{\bar{N}_\chi}^{(\delta N)}$ given in Eq. (3.2). Thus, the total angular power spectrum crucially depends on three parameters, $\kappa_b$, $\Psi^{(\delta \chi)}$, and $S_{\bar{N}_\chi}^{(\delta N)}$. To parameterize the relative size between $\Psi_{\text{RD}}^{(\delta \chi)}$ and $S_{\bar{N}_\chi}^{(\delta N)}$, we define

$$R_b = S_{\bar{N}_\chi}^{(\delta N)} / \Psi_{\text{RD}}^{(\delta \chi)}.$$ \hspace{1cm} (4.4)
Figure 1: The CMB angular power spectrum from the primordial fluctuation of the right-handed sneutrino. The $\kappa_b$ parameter is taken as 0, $-3$, $-10$, $-30$, $-100$, and $-\infty$ from above. We consider the flat universe with $\Omega_b h^2 = 0.02$, $\Omega_m = 0.3$, and $h = 0.65$, and the initial power spectral indices for primordial density perturbations are all assumed to be 1 (i.e., we adopt scale-invariant initial power spectra). The overall normalizations are taken as $[l(l+1)C_l^{(\delta \tilde{N})}/2\pi]_{l=10} = 1$.

(Notice that, with the above definition, $\Psi_{RD}^{(\delta \tilde{N})} = -\frac{2}{3} f_{\gamma N} R_b \Psi_{RD}^{(\delta \chi)}$. The $R_b$-parameter depends on the initial amplitude of the right-handed sneutrino as well as the model of inflation. For the chaotic inflation model, for example, $R_b \simeq 0.6 M_*/\tilde{N}_{init}$.)

Then, the shape of the total angular power spectrum is determined once the parameters $R_b$ and $\kappa_b$ (as well as other cosmological parameters) are fixed. As these parameters vary, the shape of the angular power spectrum changes as follows. As the $R_b$-parameter is increased, $C_l^{(\delta \tilde{N})}$ is more enhanced and hence the height of the acoustic peaks are suppressed relative to the SW tail. In addition, as $|\kappa_b|$ is increased, $C_l$ at high multipole is suppressed relative to that at low one since the acoustic peaks of $C_l^{(\delta \tilde{N})}$ is suppressed in this case. In Fig. 2, we plot the ratios $C_{10}/C_{2nd}$ and $C_{1st}/C_{2nd}$ on the $R_b$ vs. $\kappa_b$ plane.

As one can read off from Fig. 2, the shape of the angular power spectrum deviates from that from purely adiabatic density perturbations if $R_b$ and $\kappa_b$ are non-vanishing. Since the current data of the CMB angular power spectrum is well consistent with $C_l$ calculated from the purely adiabatic density perturbations, $C_l$ in our scenario becomes inconsistent with the experiments if the $R_b$ and $|\kappa_b|$ become too large. In order to derive the constraint on these parameters, we calculate the goodness-of-fit parameter $\chi^2 = -2 \ln L$, where $L$ is the likelihood function, as a function of $R_b$ and $\kappa_b$. In our calculation, the offset log-normal
Figure 2: The ratios $C_{10}/C_{2nd}$ (top) and $C_{1st}/C_{2nd}$ (bottom) on the $R_b$ vs. $\kappa_b$ plane. (The numbers in the figure are the corresponding ratios.) The shaded region corresponds to the parameter space with $\chi^2 > 84$. The cosmological parameters are the same as those used in Fig. 1.
approximation is used [27], and we use a data set consisting of 65 data points; 24 from COBE [16], 19 from BOOMERanG [17], 13 from MAXIMA [18], and 9 from DASI [19]. Then, in Fig. 2, we shaded the region where $\chi > 84$, which corresponds to 95 % C.L. excluded region for the $\chi^2$ statistics with 64 degrees of freedom. As expected, the angular power spectrum at lower multipoles is enhanced relative to that at higher multipoles. In particular, the the ratio $C_{10}/C_{2nd} \simeq 0.25$ in the purely adiabatic case, and hence the height of the SW tail can be enhanced by factor $\sim 2$ in the case of the sneutrino leptogenesis. In addition, $C_{1st}/C_{2nd} \simeq 2.0$ in the purely adiabatic case, and hence the enhancement of the ratio $C_{1st}/C_{2nd}$ is 5 % or so in this case.

Precise determination of the CMB angular power spectrum by the MAP experiment will provide stronger constraints on our scenario. At the MAP experiment, uncertainty in $C_l$ for multipoles $l \lesssim 1000$ is expected to be dominated by the cosmic variance, and in this case the error of $C_l$ is given by

$$\delta C_l = \sqrt{\frac{2}{2l+1} C_l}.$$ (4.5)

Thus, error of single $C_l$ may not be small. However, combining the informations derived from $C_l$ with different $l$, uncertainties can be reduced. For example, if we use the data for $2 \leq l \leq 50$, $201 \leq l \leq 250$, and $526 \leq l \leq 575$ to estimate the heights of the SW tail (represented by $C_{10}$) and the first and acoustic peaks, the errors are estimated as $\delta C_{10}/C_{10} \simeq 1.5 \%$, $\delta C_{1st}/C_{1st} \simeq 0.9 \%$, and $\delta C_{2nd}/C_{2nd} \simeq 0.6 \%$. With this accuracy, effects of isocurvature fluctuations will be observed if the isocurvature contribution enhances the height of the SW tail by a few % or so.

One important question is whether we can distinguish the case with correlated mixture of the adiabatic and isocurvature fluctuations from the case with uncorrelated isocurvature fluctuation (i.e., $\kappa_b \rightarrow \infty$). Since the angular power spectrum at low multipoles are suppressed relative to $C_l$ at higher multipoles in both cases, two cases are indistinguishable only by determining the ratio $C_{10}/C_{2nd}$ or $C_{1st}/C_{2nd}$. As seen in Fig. 2, however, contours of constant $C_{10}/C_{2nd}$ and those of $C_{1st}/C_{2nd}$ on the $R_b$ vs. $\kappa_b$ plane are not parallel and hence these two cases can be, in principle, distinguished by simultaneously determining the ratios $C_{10}/C_{2nd}$ and $C_{1st}/C_{2nd}$.

To discuss this issue, in Fig. 3, we plot the CMB angular power spectrum for $(R_b, \kappa_b) = (100, -3)$ (i.e., for the case with the correlated mixture of the adiabatic and isocurvature fluctuations) and $(6.4, -100)$ (i.e., for the case with uncorrelated isocurvature fluctuations). Here, we choose parameters such that the ratio $C_{10}/C_{2nd}$ becomes the same for two cases. Even if the ratio $C_{10}/C_{2nd}$ does not differ, the heights of the first acoustic peak are different; in this case, the ratio $C_{1st}/C_{2nd}$ differs about 5 %, which is within the reach of the MAP experiment if the error in $C_l$ mentioned before is realized. Thus, if the effect of the isocurvature mode is relatively large, an evidence of the correlated mixture of the adiabatic and isocurvature fluctuations may be observed by the MAP experiment.
Figure 3: The total CMB angular power spectrum with \((R_b, \kappa_b) = (100, -3)\) (dashed) and \((6.4, -100)\) (solid). The cosmological parameters are the same as those used in Fig. 1, and the normalizations are arbitrary. For comparison, we also plot the result in the purely adiabatic case (dot-dashed).

5 Summary

In this paper, we discussed the CMB anisotropy in the scenario where the baryon asymmetry of the universe originates in a scalar field condensation. We have seen that, in such a scenario, correlated mixture of the adiabatic and isocurvature fluctuations may be generated in particular when the decay product of the inflaton field and that of the scalar field (i.e., \(\tilde{N}\) in our example) both significantly contribute to the present CMB radiation. With such a correlated mixture of fluctuations, the CMB angular power spectrum may be significantly affected and the on-going MAP experiment may observe the signal of the baryogenesis from the scalar-field condensation.

In our discussion, we used the sneutrino-induced leptogenesis as an example to make our discussion clearer. However, our discussion can be applied to a wide class of scenarios where the baryon asymmetry of the universe originates in a scalar-field condensation, like Affleck-Dine scenario.

Acknowledgment: We acknowledge the use of CMBFAST [29] and RADPACK [30] packages for our numerical calculations. T.M. would like to thank the theory group of the Lawrence Berkeley National Laboratory where this work was initiated. The work of T.M. is supported by the Grant-in-Aid for Scientific Research from the Ministry of Education, Science, Sports, and Culture of Japan, No. 12047201 and No. 13740138. The work of H.M. is supported in part by the DOE Contract DE-AC03-76SF00098 and in part by the NSF.
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