Minimal supergravity, inflation, and all that

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Abstract

We consider an inflationary model in the hidden-sector broken supergravity with an effectively large cutoff. The inflaton decay into right-handed neutrinos naturally causes the observed baryon asymmetry of the universe with a reheating temperature low enough to avoid the gravitino overproduction. We emphasize that all the phenomenological requirements from cosmology and particle physics are satisfied in the large-cutoff theory.
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1. Introduction

The landscape of many vacua[1] is a plausible structure in the fundamental theory of physical laws in nature. In particular, this structure is expected as one of the theoretical ingredients to understand the observed small cosmological constant [1]. However, the anthropically allowed region of vacua in the landscape seems too large to be predictive enough in the presence of a variety of couplings. Thus, it is a challenging problem to derive further physical consequences from the landscape of vacua.

The (non-)presence of inflationary dynamics is a promising candidate as the first criterion to select realistic vacua [2]. We can naturally expect that macroscopic universe is realized through inflation from fundamental-scale physics. Moreover, under the dynamics of inflation, mediocrity principle [3] may prefer long-lasting inflations which result in larger-volume universes where more habitable galaxies are produced. In this respect, multiple inflations [4] give a remarkable possibility to be considered [5].

In a recent article [5], we have pointed out that the inflationary dynamics possess a potential to select minimal supergravity

as a large-cutoff theory, where the gravitational scale M_G is smaller than the cutoff scale M_* stemming from the fundamental theory. Such a large-cutoff supergravity naturally causes multiple slow-roll inflations, which possibly meet mediocrity principle.

The large-cutoff theory is also attractive from the viewpoint of particle-physics phenomenology: First of all, the suppression of the flavor-changing neutral currents is automatic in the large-cutoff theory, since all of the higher-dimensional operators are suppressed by the large cutoff M_* except for the genuine gravitational interactions. In detail, the large-cutoff supergravity predicts a hierarchical spectrum [5] of supersymmetric particles as $m_0 \gg |M_i|$, where m_0 is the universal soft mass for sfermions and M_i the gaugino masses ($i = 1, 2, 3$). Thus, the current chargino mass bound suggests heavy sfermions at several TeV. Such a soft mass parameter belongs to the parabolic [5] or hyperbolic [6] regime allowed for a given μ parameter. Indeed the recent detailed analysis [5] has confirmed that the region with large sfermion masses along the small-μ-parameter curve continued from the focus point [7] is consistent with the electroweak symmetry breaking [8]. In the region of heavy sfermions and light gauginos (an order of magnitude lighter than the sfermions), the constraint from CP violation is rather weak and even order one CP-violating phases are allowed for $m_0 \gtrsim 10$ TeV [9]. Furthermore, the lightest supersymmetric particle can explain the dark matter density of the present

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[1] The vacua here have extended meaning which indicates the backgrounds in the theory (moduli) space, or the landscape.

universe in the above mass region for supersymmetric particles [10].

In this Letter, we discuss a minimal new inflation model[2] as an example in the framework of the large-cutoff supergravity with emphasis on baryon asymmetry generated by leptogenesis [11] to complete a model of the large-cutoff hypothesis.[3] We claim that all the phenomenological requirements from cosmology and particle physics are satisfied in a certain parameter region of the large-cutoff theory.

2. Supergravity new inflation

We adopt a new inflation model considered in Refs. [13,14]. As an effective field theory for an inflaton chiral superfield $\tilde{\phi}$, the superpotential is given by

$$W = \tilde{v}^2 \tilde{\phi} - \frac{\tilde{g}}{n+1} \tilde{\phi}^{n+1}, \tag{1}$$

for $n \geqslant 3$ and the Kähler potential is given by

$$K = \tilde{K}|\tilde{\phi}|^2 + \frac{\tilde{k}}{4}|\tilde{\phi}|^4 + \cdots, \tag{2}$$

where we have taken the unit with the reduced Planck scale $M_G \simeq 2.4 \times 10^{18}$ GeV equal to one. The positive parameters $\tilde{K}$, $\tilde{g}$, and $\tilde{k}$ are of orders 1, $10^{-(n-2)}$, and 10^{-2}, respectively, for our large-cutoff hypothesis $M_* \simeq 10 M_G$. The tiny scale $\tilde{v}^2 > 0$ can be generated dynamically [15] and the ellipsis denotes higher-dimensional operators which may be neglected in the following analysis.

For the canonically normalized field $\phi = \sqrt{\tilde{K}}\tilde{\phi}$, the superpotential is given by [16]

$$W = v^2 \phi - \frac{g}{n+1} \phi^{n+1}, \tag{3}$$

and the Kähler potential is given by

$$K = |\phi|^2 + \frac{k}{4}|\phi|^4 + \cdots, \tag{4}$$

where we have defined

$$\tilde{v}^2 = v^2 \sqrt{\tilde{K}}, \qquad \tilde{g} = g \tilde{K}^{\frac{n+1}{2}}, \qquad \tilde{k} = k \tilde{K}^2. \tag{5}$$

The effective potential for the lowest component of ϕ is given by

$$V = e^K \left\{ \left(\frac{\partial^2 K}{\partial \phi \partial \phi^\dagger} \right)^{-1} |DW|^2 - 3|W|^2 \right\}, \tag{6}$$

where

$$DW = \frac{\partial W}{\partial \phi} + \frac{\partial K}{\partial \phi} W. \tag{7}$$

Thus, the potential of the inflaton field $\varphi = \sqrt{2}\,\mathrm{Re}\,\phi$ is approximately given by

$$V(\varphi) \simeq v^4 - \frac{k}{2} v^4 \varphi^2 - \frac{g}{2^{\frac{n}{2}-1}} v^2 \varphi^n + \frac{g^2}{2^n} \varphi^{2n} \tag{8}$$

for the inflationary period near the origin $\varphi > 0$.

The inflationary regime is determined by the slow-roll condition [17]

$$\epsilon(\varphi) = \frac{1}{2} \left(\frac{V'(\varphi)}{V(\varphi)} \right)^2 \leqslant 1, \qquad |\eta(\varphi)| \leqslant 1, \tag{9}$$

where

$$\eta(\varphi) = \frac{V''(\varphi)}{V(\varphi)}. \tag{10}$$

For the potential Eq. (8), we obtain

$$\epsilon(\varphi) \simeq \frac{1}{2} \left(\frac{-kv^4\varphi - \frac{gn}{2^{\frac{n}{2}-1}} v^2 \varphi^{n-1}}{v^4} \right)^2, \tag{11}$$

$$\eta(\varphi) \simeq \frac{-kv^4 - \frac{g}{2^{\frac{n}{2}-1}} n(n-1) v^2 \varphi^{n-2}}{v^4}. \tag{12}$$

The slow-roll condition Eq. (9) is satisfied for $\varphi \leqslant \varphi_f$ where

$$\varphi_f \simeq \sqrt{2} \left(\frac{(1-k)v^2}{gn(n-1)} \right)^{\frac{1}{n-2}}, \tag{13}$$

which yields the value of the inflaton field at the end of inflation.

The value φ_{N_e} of the inflaton corresponding to the e-fold number N_e is given by

$$N_e \simeq \int_{\varphi_f}^{\varphi_{N_e}} d\varphi \frac{V(\varphi)}{V'(\varphi)} \simeq \int_{\varphi_f}^{\varphi_{N_e}} d\varphi \frac{v^4}{-kv^4\varphi - \frac{gn}{2^{\frac{n}{2}-1}} v^2 \varphi^{n-1}}. \tag{14}$$

This leads to

$$\varphi_{N_e}^{n-2} \simeq \frac{kv^2 2^{\frac{n}{2}-1}}{gn} \left\{ \frac{1+k(n-2)}{1-k} e^{N_e k(n-2)} - 1 \right\}^{-1}. \tag{15}$$

Hence the spectral index of the density fluctuations is given by [14]

$$n_s \simeq 1 - 6\epsilon(\varphi_{N_e}) + 2\eta(\varphi_{N_e}) \tag{16}$$

$$\simeq 1 - 2k \left[1 + \frac{n-1}{\{1 + \frac{k}{1-k}(n-1)\}e^{N_e k(n-2)} - 1} \right]. \tag{17}$$

Note that the spectral index does not depend on v^2 and g explicitly. We show the k dependence of the spectral index n_s for $n = 4$ and $N_e = 45, 50, 55$ [13,14] in Fig. 1.

Now we proceed to determine the inflation scale v from the density fluctuations. The amplitude of primordial density fluctuations is given by

$$\frac{\delta\rho}{\rho} \simeq \frac{1}{5\sqrt{3}\pi} \frac{V^{\frac{3}{2}}(\varphi_{N_0})}{|V'(\varphi_{N_0})|} \simeq \frac{1}{5\sqrt{3}\pi} \frac{v^6}{kv^4\varphi_{N_0} + \frac{gnv^2}{2^{\frac{n}{2}-1}} \varphi_{N_0}^{n-1}}, \tag{18}$$
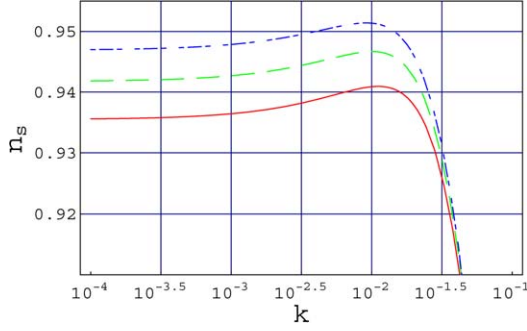
Fig. 1. (Colour online.) The k dependence of the spectral index n_s for $n = 4$. The red (solid) line corresponds to the e-fold number $N_e = 45$, the green (dashed) line to $N_e = 50$, and the blue (dash-dotted) line to $N_e = 55$. For $k = 0$, $n_s \simeq 1 - 6/(2N_e + 3)$.
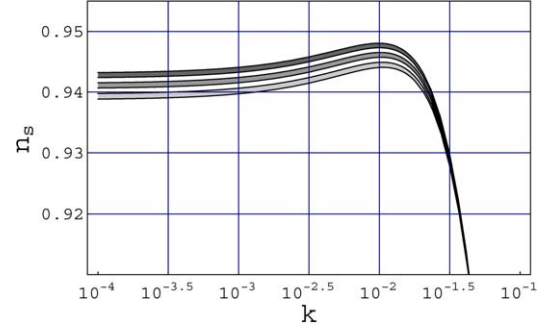

Fig. 2. The k dependence of the spectral index n_s for $n = 4$. The shaded regions correspond to $T_R = 10^5, 10^7, 10^9$ GeV from below, and the lower lines for $g = 0.1$ and the upper lines for $g = 0.01$.

where φ_{N_0} is the value of inflaton field at the epoch of the present-horizon exit. Thus we obtain

$$v^{\frac{2n-6}{n-2}} \simeq \sqrt{2}\, \frac{V^{\frac{3}{2}}(\varphi_{N_0})}{|V'(\varphi_{N_0})|} \left[\frac{k}{gn} \left\{ \frac{1 + k(n-2)}{1-k} e^{N_0 k(n-2)} - 1 \right\}^{-1} \right]^{\frac{1}{n-2}}$$
$$\times \left[k + k \left\{ \frac{1 + k(n-2)}{1-k} e^{N_0 k(n-2)} - 1 \right\}^{-1} \right]. \tag{19}$$

Owing to the COBE normalization

$$\frac{V^{\frac{3}{2}}(\varphi_{N_0})}{|V'(\varphi_{N_0})|} \simeq 5.3 \times 10^{-4}, \tag{20}$$

the scale v is expressed as

$$v \simeq 10^{12}\ \text{GeV} \times \mathcal{C}(k, N_0) \times \left(\frac{0.1}{g} \right)^{\frac{1}{2}}, \tag{21}$$

for $n = 4$, $N_0 \simeq 50$, and $k \lesssim 0.01$, where $\mathcal{C}(k, N_0)$ is a function of order unity.

On the other hand, the e-fold number of the present horizon is also given by

$$N_0 \simeq 67 + \frac{1}{3} \ln H + \frac{1}{3} \ln T_R$$
$$\simeq 67 + \frac{1}{3} \ln \frac{v^2}{\sqrt{3}} + \frac{1}{3} \ln T_R, \tag{22}$$

where H denotes the Hubble scale at the horizon exit and T_R the reheating temperature. By means of Eqs. (17), (19), and (22), we can determine v and N_0 from g, k, and T_R. For $n = 4$, $g \sim 0.1$, $k \sim 0.01$, and $T_R \sim 10^{5-9}$ GeV, the inflation scale v is given by $\mathcal{O}(10^{12})$ GeV, and the e-fold number N_0 of the present horizon is given by 47.6–50.6. In Fig. 2, we show the k dependence of the spectral index n_s for the reheating temperature $T_R = 10^5$, 10^7, and 10^9 GeV. We conclude that the implication $\tilde{g} \sim \tilde{k} \sim 0.01$ of the large-cutoff hypothesis[4] is consistent with an experimental value $n_s = 0.95 \pm 0.02$ [18] of the spectral index for a wide range of the reheating temperature.[5]

3. The gravitino mass

In the previous section, we have confirmed that the new inflation model in the large-cutoff hypothesis is consistent with the cosmological observations. In this section, we discuss the gravitino problem under such an inflationary scenario.

As considered in Ref. [13], we assume that the positive energy Λ^4_{SUSY} of the SUSY breaking is dominantly canceled out by the negative energy at the inflaton potential minimum. Namely we impose

$$\Lambda^4_{\text{SUSY}} - 3|W(\phi_0)|^2 = 0, \tag{23}$$

where ϕ_0 is the minimum point of ϕ in Eq. (6).

Then we obtain the gravitino mass as

$$m_{\frac{3}{2}} \simeq \frac{\Lambda^2_{\text{SUSY}}}{\sqrt{3}} = W(\phi_0). \tag{24}$$

The value of ϕ_0 is approximately given by

$$\phi_0 \simeq \left(\frac{v^2}{g} \right)^{\frac{1}{n}}. \tag{25}$$

Consequently the gravitino mass is given by

$$m_{\frac{3}{2}} \simeq \frac{n v^2}{n+1} \left(\frac{v^2}{g} \right)^{\frac{1}{n}} \simeq 9\ \text{TeV} \times \left(\frac{0.1}{g} \right)^{\frac{3}{2}}. \tag{26}$$

The second equality holds for $n = 4$, where we have used Eq. (21) and omitted the weak dependence on k and T_R.

More precisely, by means of Eqs. (19) and (22), the gravitino mass can be expressed as a function of g, k, and T_R, although the dependence on T_R is very weak, as can be seen

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amplitude $\Delta\varphi$ at $\varphi = \varphi_{N_0}$ is given by

$$\Delta\varphi|_{\varphi_{N_0}} \sim \frac{\sqrt{2\epsilon(\varphi_{N_0})}}{2\pi\sqrt{3}} \frac{V(\varphi_{N_0})^{\frac{3}{2}}}{V'(\varphi_{N_0})}$$
$$\simeq \frac{1}{2\pi\sqrt{3}} \frac{V(\varphi_{N_0})^{\frac{3}{2}}}{V'(\varphi_{N_0})} \left(k + \frac{gn}{2^{\frac{n}{2}-1}} \frac{\varphi_{N_0}^{n-2}}{v^2} \right) \varphi_{N_0}.$$

For $n = 4$, $g \sim 0.1$, $k \sim 0.01$, and $T_R \sim 10^{5-9}$ GeV, the fluctuation amplitude $\Delta\varphi|_{\varphi_{N_0}}$ takes a value of order $10^{-6}\varphi_{N_0}$, which is much less than the mean-field value φ_{N_0} to justify the above slow-roll analysis.
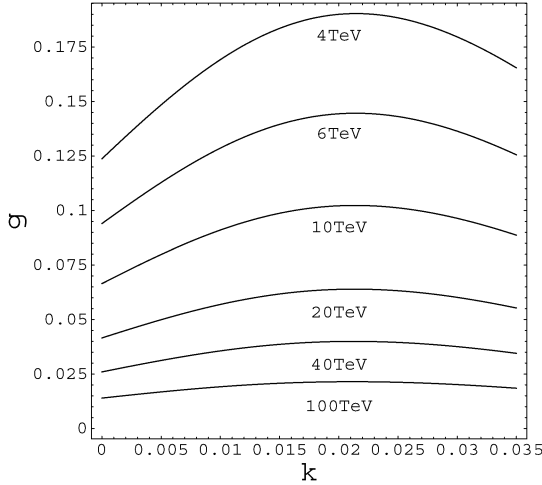
Fig. 3. The contours of the gravitino mass for $n = 4$ and $T_R = 4 \times 10^6$ GeV. The dependence on the reheating temperature is very weak.

from Eqs. (19) and (22). The result is shown in Fig. 3. For $g \lesssim 0.2$ and $k \lesssim 0.035$, the gravitino mass is larger than 4 TeV, which may avoid the gravitino overproduction for a reheating temperature $T_R \sim 10^{6-7}$ GeV [19].

In contrast, the sfermion soft mass is given as $m_0 \simeq m_{3/2}$ if no D-term contributes to the SUSY breaking. Thus, $m_0 < 10$ TeV implies $g \gtrsim 0.07$ for $k \lesssim 0.035$.

4. Reheating for baryogenesis

Now we are ready to consider the baryon asymmetry in the present new inflation model with the large cutoff.

We assume the baryon asymmetry is generated by leptogenesis [11] through non-thermal production of right-handed neutrinos, as investigated in Refs. [10,11], which provides a numerical estimate

$$\frac{n_B}{s} \simeq 8.2 \times 10^{-11} \left(\frac{T_R}{10^6 \text{ GeV}} \right) \left(\frac{2m_N}{m_\phi} \right) \left(\frac{m_{\nu_3}}{0.05 \text{ eV}} \right) \frac{1}{\sin^2 \beta} \delta_{\text{eff}}. \tag{27}$$

Here m_N, m_ϕ, and m_{ν_3} are the masses of the right-handed neutrino N, the inflaton ϕ and the heaviest (active) neutrino, respectively. The phase δ_{eff} is the effective CP phase defined in Ref. [11] and $\tan \beta$ is the ratio of the vacuum expectation value of up- and down-type Higgs bosons in the MSSM. The reheating temperature is given by

$$T_R \simeq \left(\frac{10}{g_* \pi^2} \Gamma_\phi^2 \right)^{\frac{1}{4}}, \tag{28}$$

where Γ_ϕ is the decay width of the inflaton and g_* is the effective number of massless degrees of freedom to be taken as $g_* = 228.75$ numerically. Note that the inflaton mass

$$m_\phi \simeq n v^2 \left(\frac{v^2}{g} \right)^{-\frac{1}{n}} \tag{29}$$

in our new inflation model also weakly depends on the k and the reheating temperature T_R, as is the case for the gravitino mass in Eq. (26).

Let us introduce the following superpotential interaction as the dominant source of the N production[6]:

$$\delta W = \frac{h}{2(n-1)} \phi^{n-1} N^2, \tag{30}$$

where h is a positive parameter of the order of the inflaton self-coupling g.[7] The coupling Eq. (30) gives a decay width

$$\Gamma_\phi \simeq \frac{|h|^2}{16\pi} \phi_0^{2(n-2)} m_\phi. \tag{31}$$

From this decay width the reheating temperature after inflation for $n = 4$ is given by[8]

$$T_R \simeq 2.6 \times 10^6 \text{ GeV} \left(\frac{h}{0.1} \right) \left(\frac{0.1}{g} \right)^{\frac{5}{4}}, \tag{32}$$

where we have omitted the weak dependence on k in Eq. (21). Therefore, the reheating temperature $T_R \sim 10^{6-7}$ GeV is typical in this model. As mentioned above, this reheating temperature is low enough to avoid the gravitino overproduction.

Note that the operator Eq. (30) also gives the Majorana mass to the neutrino:

$$m_N = \frac{h}{n-1} \phi_0^{n-1} \simeq \frac{h}{n-1} \left(\frac{v^2}{g} \right)^{1-\frac{1}{n}}. \tag{33}$$

Thus the mass inequality $2m_N < m_\phi$, namely,

$$h < \frac{n(n-1)}{2} g, \tag{34}$$

is satisfied with a typical parameter set $g \sim h$. This is appropriate for the non-thermal production of neutrinos which leads to the non-thermal leptogenesis.

Based on the above setup, we now estimate the baryon asymmetry due to the decay of inflaton[9] in our model as a function of the couplings g, k, and the reheating temperature T_R. The baryon asymmetry is determined by four independent parameters g, k, v, and h. In terms of the observed density fluctuations, we can represent v with the other parameters. We further use the reheating temperature as an input parameter instead of h by means of Eqs. (28), (29) and (31):

$$h \simeq \sqrt{\frac{16\pi}{m_\phi}} \left(\frac{g_* \pi^2}{10 M_G^2} \right)^{\frac{1}{4}} \left(\frac{g}{v^2} \right)^{\frac{n-2}{n}} T_R. \tag{35}$$

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[6] The inflaton field is also expected to decay through couplings with light fields ψ_i in the Kähler potential such as $\sum_i c_i |\phi|^2 |\psi_i|^2$. However, the decay width $\Gamma \sim \sum_i |c_i|^2 \phi_0^2 m_\phi^3$ through these couplings is so small that we neglect such contributions.

[7] Here, we assign the same charge for ϕ and N under Z_{2n} R-symmetry, while we assign the matter parity $+$ for ϕ and $-$ for N. Hence we expect the presence of such operators as $\phi^{n-3} N^4$ in addition to Eq. (30). We do not include such operators since the operator Eq. (30) with the smallest number of N dominates the reheating and leptogenesis.

[8] The cross term between $\phi^{n-1} N^2$ and $v^2 \phi$ in the superpotential gives a comparable decay width. We neglect this contribution since it does not essentially affect our conclusions.

[9] In our setup, we also have an additional contribution to the baryon asymmetry and the gravitino abundance. However, as we see in Appendix A, this contribution is small in typical parameter region so that we neglect this contribution in the following analysis.
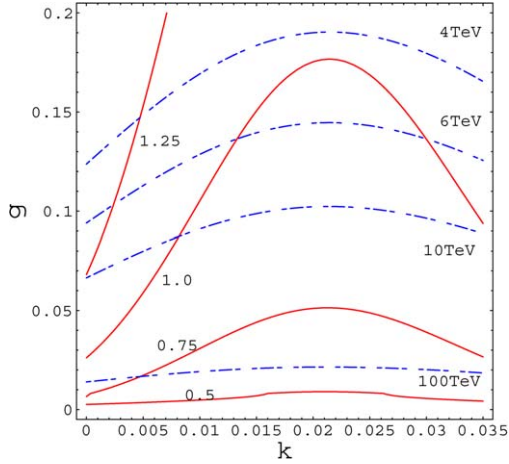
Fig. 4. (Colour online.) The contours of $(n_B/s)/(n_B/s)_0$ for $n = 4$, $T_R = 4 \times 10^6$ GeV, $\delta_{\text{eff}} = 1$, $\sin\beta = 1$ are plotted in red (solid) lines. The blue (dashed-dotted) lines correspond to the contours of gravitino mass.

Then the baryon asymmetry n_B/s is given in terms of g, k, and T_R by

$$\frac{n_B}{s} \simeq 8.2 \times 10^{-11} \left(\frac{T_R}{10^6 \text{ GeV}} \right) \left(\frac{2h}{n(n-1)g} \right)$$
$$\times \left(\frac{m_{\nu_3}}{0.05 \text{ eV}} \right) \frac{1}{\sin^2\beta} \delta_{\text{eff}}, \qquad (36)$$

where $h(g, k, T_R)$ is given by Eq. (35) with v determined by Eqs. (19) and (22).

In Fig. 4, we plot the contours of $m_{3/2}$ and $(n_B/s)/(n_B/s)_0$ for $T_R = 4 \times 10^6$ GeV, $m_{\nu_3} = 0.05$ eV, $\delta_{\text{eff}} = 1$, $\sin\beta = 1$, where $(n_B/s)_0$ is the baryon asymmetry of the universe suggested by WMAP [21]:

$$\left(\frac{n_B}{s} \right)_0 \simeq 8.7 \times 10^{-11}. \qquad (37)$$

We note that the baryon asymmetry and the gravitino mass for different reheating temperatures can also be seen from Fig. 4: The baryon asymmetry is proportional to the square of the reheating temperature T_R^2, since the coupling h is approximately proportional to the reheating temperature. As for the gravitino mass, its value is almost independent of T_R, since v is almost independent of T_R.

This figure shows that the sufficient baryon asymmetry is produced in a typical parameter region of the large-cutoff hypothesis: $k \sim 0.01$, $g \sim 0.01$–0.1, and $T_R \sim 10^6$ GeV, which turns out to be low enough to avoid the gravitino overproduction. Thus it is revealed that the large-cutoff hypothesis is also consistent with the observed baryon asymmetry.

5. Discussion

We have studied the large-cutoff hypothesis from the viewpoint of cosmology. We first confirmed that the spectral index in the new inflation model has an upper bound $n_s \lesssim 0.95$ (see Ref. [14]) and the large-cutoff hypothesis implies its boundary value, which remarkably agrees with the present experimental suggestion $n_s = 0.95 \pm 0.02$ [18]. Secondly, we found a

concrete setup where the sufficient baryon asymmetry can be produced via non-thermal leptogenesis with the reheating temperature low enough to avoid the gravitino overproduction in a typical parameter region of large-cutoff hypothesis.

We again emphasize that the large cut-off hypothesis has several advantages from the viewpoint of particle-physics phenomenology. It solves the FCNC problem and produces the mass spectrum $m_0 \sim 10m_{1/2} \sim 10\mu$, which yields the correct electroweak symmetry breaking [5]. Furthermore, the spectrum realized in the large-cutoff hypothesis accommodates the appropriate amount of the dark matter density [10].

We also mention CP violations in the visible-sector supersymmetric standard model as a sensitive low-energy probe of the supersymmetry breaking. Phases of the theory would be limited severely if the scalar masses were to be less than the TeV scale. In contrast, for $m_0 \sim 10$ TeV, such a constraint is far milder, with the very heavy scalar masses expected to be realized in the large-cutoff hypothesis from the viewpoint of electroweak symmetry breaking and dark matter, as mentioned above.

The heavy scalar masses are remarkably consistent with the cosmological constraint, as we saw in this Letter. Thus we conclude that the large-cutoff theory with the supergravity new inflation and non-thermal leptogenesis is consistent with all the phenomenological requirements from cosmology and particle physics.

Finally we comment on other types of inflations. The presence of the large cutoff seems advantageous for other inflationary models such as hybrid inflation and chaotic inflation. In particular, large-field inflations imply the presence of a larger scale (see Ref. [22]) than the reduced Planck scale. In fact, we suspect that multiple inflations may be so generic as to include various types of inflations as their components, whose slow-roll conditions are realized by the large-cutoff mechanism [5].

Acknowledgements

Appendix A. Another source of baryon and gravitino

In Section 4, we put aside the baryon asymmetry and the gravitino produced through the coherent oscillation of right-handed sneutrino. Here we argue that this contribution can be small enough to be neglected.

Firstly we explain the motion of right-handed sneutrino field which is the source of the baryon asymmetry and gravitino. During inflation, the right-handed sneutrino is fixed at the origin due to the Hubble mass. After the inflaton starts to roll down to the vacuum, the mass of the right-handed sneutrino changes along the motion trajectory of the inflaton. As the oscillation energy of the inflaton decreases, the origin of right-handed sneutrino becomes unstable, and right-handed sneutrino also starts oscillation. Then the decay of right-handed sneutrino becomes

significant.[10] The baryon asymmetry and gravitino are provided by the decay of this right-handed sneutrino [20].

Let us estimate the yields of the baryon asymmetry and gravitino provided through the coherent oscillation of right-handed sneutrino. As mentioned above, the decay of right-handed sneutrino becomes significant when the motion of right-handed sneutrino is induced by that of inflaton. Then the yields of the baryon asymmetry n_B^N/s and the gravitino number $n_{3/2}^N/s$ produced at the decay time of right-handed sneutrino are given by

$$\frac{n_B^N}{s} \simeq \varepsilon \frac{\rho_N}{m_N} \frac{45}{2\pi^2 g_* T_N^3}, \tag{A.1}$$

$$\frac{n_{3/2}^N}{s} \simeq Y_{3/2}^{\phi} \frac{T_N}{T_R}. \tag{A.2}$$

Here, ε denotes the CP-asymmetry in right-handed sneutrino decay defined in Ref. [11], T_N is the temperature of radiation produced by right-handed sneutrino decay, $Y_{3/2}^{\phi}$ is the yield of gravitino produced by inflaton decay, and ρ_N is the energy of the right-handed sneutrino at the right-handed sneutrino decay.

After the inflaton decay, these yields are diluted by the dilution factor Δ estimated as

$$\Delta \simeq \frac{T_N}{T_R} \frac{\rho_\phi}{\rho_N}, \tag{A.3}$$

where ρ_ϕ is the energy of the inflaton at the right-handed sneutrino decay. Thus n_B^N/s and $n_{3/2}^N/s$ after the inflaton decay are given by

$$\frac{n_B^N}{s} \simeq \varepsilon \frac{\rho_N}{m_N} \frac{45}{2\pi^2 g_* T_N^3} \frac{T_R}{T_N} \frac{\rho_N}{\rho_\phi}$$

$$\simeq 5.3 \times 10^{-11} \left(\frac{T_R}{10^6 \text{ GeV}}\right) \left(\frac{m_{\nu_3}}{0.05 \text{ eV}}\right) \frac{\rho_N}{\rho_\phi} \delta_{\text{eff}}, \tag{A.4}$$

$$\frac{n_{3/2}^N}{s} \simeq Y_{3/2}^{\phi} \frac{\rho_N}{\rho_\phi}. \tag{A.5}$$

These values are smaller than the yields produced at inflaton decay for $\rho_N \ll \rho_\phi$ (see Eq. (36)), which we assume in the main text.

In fact, we checked that $\rho_N \ll \rho_\phi$ is realized in a typical parameter region by solving the equations of motion numerically for $n = 4$. We note a possibility that parametric resonance occurs in specific points, and the energy of right-handed sneutrino ρ_N becomes comparable to that of inflaton ρ_ϕ in such a case.

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[10] The decay width of right-handed sneutrino is much larger than that of inflaton, due to a large Yukawa coupling of right-handed neutrino and standard-model particles compared with Eq. (30).

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