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Spontaneous Baryogenesis from Asymmetric Inflaton

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We propose a variant scenario of spontaneous baryogenesis from asymmetric inflaton based on current-current interactions between the inflaton and matter fields with a non-zero $B - L$ charge. When the inflaton starts to oscillate around the minimum after inflation, it may lead to excitation of a CP-odd component, which induces an effective chemical potential for the $B - L$ number through the current-current interactions. We study concrete inflation models and show that the spontaneous baryogenesis scenario can be naturally implemented in the chaotic inflation in supergravity.

I. INTRODUCTION

The history of the Universe after the Big Bang Nucleosynthesis epoch is well understood, and it is known as the standard Big Bang cosmology. However, it suffers from various initial condition problems such as the horizon problem, flatness problem, and the origin of density fluctuations. In the inflationary paradigm [1–5], the exponential expansion of the Universe solves these problems. On the other hand, in the standard Big Bang cosmology one needs to adopt an initial condition such that the amount of baryon number is ten orders of magnitude smaller than the entropy density. In the inflationary Universe, any pre-existing baryon asymmetry would be exponentially diluted, and so, the baryon asymmetry needs to be created after inflation.

The inflation must be connected to the subsequent hot Big Bang phase. This is naturally realized in the slow-roll inflationary scenario [6, 7], where the inflation is driven by a scalar field called inflaton which slowly rolls down the nearly flat potential. While the Universe is dominated by the potential energy of the inflaton, it experiences an exponential expansion. The inflation ends when the inflaton starts to oscillate around the potential minimum, and the potential energy is converted to radiation through the inflaton decay.

Suppose that the inflaton is a part of a complex scalar field with an approximate U(1) symmetry around the potential minimum. This is often the case in supersymmetric (SUSY) theories, where each chiral multiplet contains a complex scalar. When the inflaton starts to oscillate about the potential minimum after inflation, it may acquire a non-zero U(1) asymmetry associated with the inflaton number. (Here the inflaton number refers to the CP-odd component of the excited inflaton quanta.) Once produced, the inflaton number decreases as a^{-3} due to the expansion of the Universe, where a is the scale factor. If there are current-current interactions between the inflaton number and the $B-L$ symmetry, the excited CP-odd inflaton quanta induces an effective chemical potential of the $B - L$ number. This leads to the spontaneous baryogenesis if the $B - L$ number is broken in thermal plasma [8–11], because the inflaton asymmetry biases the $B - L$ number. As for the $B - L$ breaking, one may introduce dimension five interactions for the neutrino mass, motivated by the seesaw mechanism [12]. Finally the inflaton decays into radiation and reheats the Universe, which connects the inflation to the hot Big Bang phase.

In this paper, we consider the spontaneous baryogenesis scenario based on the current-current interactions between the inflaton and $B - L$ numbers. We evaluate the abundance of baryon asymmetry and clarify what conditions are needed to explain the observed amount of the baryon asymmetry. We also study explicit inflation models and show that the scenario can be naturally implemented in the chaotic inflation in supergravity. Interestingly, no isocurvature perturbations are generated in this case, in contrast to the usual spontaneous baryogenesis [13].

The rest of this paper is organized as follows. In the next section we will explain our main idea about the spontaneous baryogenesis from asymmetric inflaton. In Sec. III we study concrete inflation models in supergravity. The last section is devoted for discussion and conclusions.

II. SPONTANEOUS BARYOGENESIS FROM ASYMMETRIC INFLATON

Suppose that inflaton ϕ is a complex scalar field with an approximate conserved $\mathrm{U}(1)_{\rm inf}$ current of

$$j_\mu^{\rm inf} = 2{\rm Im}\,(\phi\partial_\mu\phi^*) + \dots . \quad (1)$$

where ... represents the other fields carrying nonzero inflaton charges such as the inflatino. This is the case in many SUSY inflation models because each chiral multiplet contains a complex scalar field. In a class of inflation models, the flatness of the inflaton potential is due to the $\mathrm{U}(1)_R$ symmetry, in which case the inflaton current is identified with the $\mathrm{U}(1)_R$ current. We assume that the $\mathrm{U}(1)_{\rm inf}$ symmetry is only an approximate symmetry at the potential minimum, but it is generically broken explicitly and in particular, it is generically spontaneously broken during inflation. When the inflaton starts to oscillate about the potential minimum after inflation, a nonzero inflaton number may be produced if the $\mathrm{U}(1)_{\rm inf}$ is explicitly broken. The inflaton number decreases as a^{-3} if the breaking term is irrelevant in the vicinity of the potential minimum. This is naturally realized in some models as shown in the next section. If the inflaton current is coupled to the $B-L$ current, the CP asymmetric inflaton number biases the $B-L$ number. Thus, spontaneous baryogenesis takes place if $B-L$ number is explicitly broken in plasma. Some standard model (SM) particles may carry a nonzero inflaton charge to realize the reheating. We consider the case that the inflaton charge operator commutes with the $B-L$ charge so that we focus on baryon asymmetry generated via the spontaneous baryogenesis. Otherwise the baryon asymmetry is directly generated by inflaton dynamics, which is the case out of our interest (see, e.g., Ref. [14] for that case).

To realize spontaneous baryogenesis, we introduce the following current-current interaction:

$$-\mathcal{L} = \tilde{G}_F j^{\rm inf}_\mu j^\mu_{B-L}, \tag{2}$$

where $\tilde{G}_F$ is an effective coupling constant with mass dimension minus two.[1] Here, j_{B-L} is the $B-L$ current,

$$j^\mu_{B-L} = \sum_i q_i j^\mu_i \tag{3}$$

$$j^\mu_i = \begin{cases} \bar{\psi}_i \gamma^\mu \psi_i & \text{for fermions} \\ 2\mathrm{Im}\left(\varphi_i \partial^\mu \varphi_i^*\right) & \text{for bosons,} \end{cases} \tag{4}$$

where q_i is the $B-L$ charge of the field i. Such current-current interactions are generically present in supergravity theories as shown in Sec. III.

After inflation ends, the coherent oscillations of the inflaton dominate the energy density of the Universe. The Friedmann equation implies the following relation:

$$n_{\rm inf} \simeq \frac{3\epsilon H^2(t) M_{\rm Pl}^2}{m_{\rm inf}} \tag{5}$$

$$\epsilon \equiv \frac{m_{\rm inf} n_{\rm inf}}{\rho_{\rm inf}}, \tag{6}$$

where $M_{\rm Pl}$ ($\simeq 2.4 \times 10^{18}$ GeV) is the reduced Planck mass, $m_{\rm inf}$ is the inflaton mass at the potential minimum, $n_{\rm inf} = (j^{\rm inf})^0$ is the inflaton number density and $\rho_{\rm inf}$ is the energy density of the inflaton. Here we have assumed that the only inflaton has a sizable $\mathrm{U}(1)_{\rm inf}$ asymmetry. We define ϵ (≤ 1) so that it represents the ellipticity of inflaton trajectory in its complex plane. As the inflaton number and energy densities are (nearly) spatially homogeneous, the relevant part of the current-current interaction is given by

$$-\mathcal{L} = \mu_{B-L} n_{B-L} \tag{7}$$

$$\mu_{B-L} \simeq 3\epsilon \tilde{G}_F M_{\rm Pl}^2 \frac{H^2(t)}{m_{\rm inf}}, \tag{8}$$

where $n_{B-L} = j^0_{B-L}$ is the $B-L$ number density. This means that the $B-L$ number has an effective chemical potential μ_{B-L}. The chemical potential biases the $B-L$ asymmetry in the chemical equilibrium as

$$n^{\rm (eq)}_{B-L} \simeq k \mu_{B-L} T^2 \tag{9}$$

$$k \equiv \sum_i q_i \frac{g_i}{6}, \tag{10}$$

where the summation is taken for all particles in the thermal bath and g_i is the number of spin states but with an extra factor of 2 for bosons. In the SM, the coefficient k is $1/2$, while in the MSSM it is 1.

In order to generate a non-zero $B-L$ asymmetry, we introduce the following $B-L$ violating interaction

$$\mathcal{L} = \frac{y^2}{2M_R}(L\tilde{H})^2 + \text{h.c.} \tag{11}$$

for the light neutrino masses, which is obtained after integrating out heavy right-handed neutrinos in the seesaw mechanism [12]. Here $\tilde{H} = i\sigma_2 H^*$ is the SU(2) conjugate of the SM Higgs doublet H. Throughout this paper we assume that the right-handed neutrinos are so heavy that they are not produced from thermal scattering, and so, our scenario is complementary to thermal leptogenesis [15]. The effective rate of the lepton number violating processes via the above interaction is roughly given by [16]

$$\Gamma_L \sim \sigma_R T^3 \sim \frac{\bar{m}^2 T^3}{16\pi v_{\rm ew}^4}, \tag{12}$$

where $v_{\rm ew}$($\simeq$ 246 GeV) is the Higgs vacuum expectation value (VEV) and $\bar{m}^2$ is the sum of the left-handed neutrino mass squared. We assume $\bar{m}^2$ to be of order the atmospheric neutrino mass squared difference, $\Delta m^2_{\rm atm} \simeq 2.4 \times 10^{-3}$ eV^2. (See e.g. Refs. [17, 18] for the latest combined-data fit parameters.) Thus we obtain $\sigma_R M_{\rm Pl}^2 \simeq 8 \times 10^4$.

Now let us calculate the baryon asymmetry. Before the reheating completes, the temperature of the plasma is written as

$$T \simeq \left(\frac{36 H(t) \Gamma_I M_{\rm Pl}^2}{g_* \pi^2}\right)^{1/4}, \tag{13}$$

where g_* is the effective relativistic degrees of freedom and equal to 106.75 (228.75) in the SM (MSSM). The inflaton decay rate Γ_I is related to the reheating temperature as

$$T_{\rm RH} \simeq \left(\frac{90}{g_* \pi^2}\right)^{1/4} \sqrt{\Gamma_I M_{\rm Pl}}. \tag{14}$$

Then the ratio between the reaction rate of the $B-L$ violating interaction and the Hubble expansion rate is given by

$$\frac{\Gamma_L}{H} \simeq 0.1 \left(\frac{H_{\rm RH}}{H(t)}\right)^{1/4} \left(\frac{T_{RH}}{M_{\rm Pl}}\right) M_{\rm Pl}^2 \sigma_R, \tag{15}$$

[1] In Ref. [11], a four-Fermi interaction between DM and $B-L$ currents with $\tilde{G}_F \sim 1/(10\ \mathrm{TeV})^2$ and the associated spontaneous baryogenesis were studied in an asymmetric dark matter model.

where $H_{\rm RH}(=\Gamma_I)$ is the Hubble parameter at the reheating. When this ratio is larger than unity, the $B-L$ asymmetry would reach the equilibrium value. However, as can be seen from the above expression, the ratio is almost always below unity, so that the $B-L$ number at the time of reheating is estimated by integrating the Boltzmann equation:

$$n_{B-L}|_{\rm RH} \simeq \int_{t_{\rm inf}}^{t_{\rm RH}} {\rm d}t' \frac{H_{\rm RH}^2}{H^2(t')}\Gamma_L n_{B-L}^{\rm (eq)}(t'), \qquad (16)$$

where $t_{\rm RH}$ and $t_{\rm inf}$ represent the cosmic time at the reheating and the end of inflation, respectively. Since the integrand is proportional to T^5 (i.e., $\propto t^{-5/4}$), the baryon abundance is mostly generated just after the end of inflation. This is in contrast with the ordinary scenario of the spontaneous baryogenesis, where the baryon abundance is mostly generated at the time of reheating. This is because in our case the effective chemical potential decreases as $\propto H^2(t)$, which is faster than the ordinary case, and the bias effect is most efficient just after the end of inflation. Thus we can estimate the $B-L$ number at the reheating as

$$n_{B-L}|_{\rm RH} \simeq \left.\frac{H_{\rm RH}^2}{H^2(t)}\frac{\Gamma_L}{H(t)} n_{B-L}^{\rm (eq)}(t)\right|_{t=t_{\rm inf}}, \qquad (17)$$

Combining those results, we obtain the final baryon asymmetry as

$$Y_b \equiv \frac{n_b}{s} \simeq \left.\frac{8n_f+4n_H}{22n_f+13n_H} T_{\rm RH}\frac{n_{B-L}}{4H^2M_{\rm Pl}^2}\right|_{\rm RH} \qquad (18)$$

$$\simeq 0.01\,\epsilon(M_{\rm Pl}^2\sigma_R)(\tilde{G}_F M_{\rm Pl}^2)\frac{H(t)^{1/4}T_{\rm RH}^{7/2}}{m_{\rm inf}M_{\rm Pl}^{11/4}}, \qquad (19)$$

where the pre-factor in the first line is due to the sphaleron effect [19]. In the SM (MSSM), $n_f=3$ and $n_H=1$ (2). In the second line, we have substituted some parameters, including g_*, n_h, and k, and obtain the factor of about 0.01 for the SM and MSSM. The observed baryon abundance of $Y_b^{\rm (obs)} \simeq 8.6\times10^{-11}$ [20] requires that the reheating temperature is as large as

$$T_{RH} \simeq 3\times10^{13}\ {\rm GeV}\epsilon^{-2/7}\left(\tilde{G}_F M_{\rm Pl}^2\right)^{-2/7} \times\left(\frac{m_{\rm inf}}{10^{13}\ {\rm GeV}}\right)^{2/7}\left(\frac{H(t)}{10^{13}\ {\rm GeV}}\right)^{-1/14}. \qquad (20)$$

Lastly let us comment on the washout effect. Around the time of reheating, the generated lepton asymmetry is partially washed out due to the inverse processes. Calculating the Boltzmann equation, we obtain a washout factor such as [10]

$$\Delta_w \simeq \exp\left[-0.7\left(\frac{T_{\rm RH}}{10^{13}\ {\rm GeV}}\right)\right]. \qquad (21)$$

This implies that the reheating temperature cannot be much larger than 10^{13} GeV to avoid the washout effect due to the inverse processes. Taking into account the washout factor, we find that the resulting baryon asymmetry has a peak around the reheating temperature of 5×10^{13}, where the result of Eq. (19) is overestimated by a factor of 5. Thus we may assume $(\tilde{G}_F M_{\rm Pl}^2)\simeq 5$ to explain the observed baryon asymmetry.

III. INFLATION MODELS

In this section, we consider concrete models of inflation in supergravity to see if the scenario in the previous section can be realized in realistic inflation models. Hereafter, we adopt the Planck units ($M_{\rm Pl}=1$).

In supergravity, the Lagrangian is constructed to be invariant in terms of supergravity transformation. The relevant part of the Lagrangian comes from kinetic terms such as

$$\mathcal{L} = K_{ij^*}\partial_\mu\varphi^i\partial^\mu\varphi^{*j} + iK_{ij^*}\bar{\chi}^j\bar{\sigma}^\mu\tilde{\mathcal{D}}_\mu\chi^i, \qquad (22)$$

$$\tilde{\mathcal{D}}_\mu\chi^i \equiv \partial_\mu\chi^i + K^{im^*}K_{km^*l}\partial_\mu\varphi^l\chi^k - \frac{1}{4}\left(K_j\partial_\mu\varphi^j - K_{j^*}\partial_\mu\varphi^{*j}\right)\chi^i + \ldots, \qquad (23)$$

where scalar and fermionic components of chiral superfields are denoted by φ_i and χ_i, respectively, and K is the Kähler potential, $K^{ij^*}\equiv(K_{ij^*})^{-1}$, and the subscripts represent derivatives with respect to the corresponding field, such as $K_{ij^*}\equiv\partial^2K/\partial\varphi^i\partial\varphi^{*j}$. Note that the total Lagrangian is real.

A. F-term hybrid inflation

Suppose that there is a Kähler potential of

$$K = |\phi|^2 + c\,|\phi|^2|\chi|^2 + |\chi|^2, \qquad (24)$$

where the scalar component of ϕ is inflaton and χ is a $B-L$ charged field. Neglecting the fermionic component of ϕ and denoting the scalar and fermionic components of χ as $\tilde{\chi}$ and χ, respectively, we obtain the desirable interaction term such as

$$\mathcal{L}_{\rm int} = {\rm Im}\left(\phi\partial_\mu\phi^*\right)\left[2c{\rm Im}\left(\tilde{\chi}\partial^\mu\tilde{\chi}^*\right) + \left(c-\frac{1}{2}\right)\bar{\chi}\bar{\sigma}^\mu\chi\right] \qquad (25)$$

where we have assumed $c\,|\phi|^2 \ll 1$ and have rescaled the fields $\tilde{\chi}$ and χ to make their kinetic term canonical. These interactions are nothing but the current-current interaction of the form (2) with $(\tilde{G}_F M_{\rm Pl}^2)\approx c$.

The above calculation can be applied to, e.g., the F-term hybrid inflation model with the superpotential [21, 22]

$$W = \lambda\phi\left(\psi\bar{\psi} - \frac{v^2}{2}\right), \qquad (26)$$

where ψ and $\bar{\psi}$ are waterfall fields, and ϕ is the inflaton with $U(1)_R$ charge 2. Note that in this case the inflaton current is nothing but the R-current. The $U(1)_R$ symmetry is necessarily broken by a constant term of the superpotential, $W_0 \simeq m_{3/2}$, which is required for realizing a vanishingly small cosmological constant. As a result, the inflaton potential receives a linear term potential in proportion to the gravitino mass, $m_{3/2}$, which modifies the inflaton dynamics [23, 24]. In particular, the angular motion of ϕ, i.e., the inflaton number, is induced during inflation [24]. The effect of the angular motion on the density perturbations was recently studied in detail in Ref. [25]. The dynamics of inflaton during inflation is mainly determined by the Coleman-Weinberg potential and the linear term of inflaton. The following parameter ξ measures the relative importance of the two contributions to the slope of the potential [25]:

$$\xi \equiv \frac{2^{9/2}\pi^2}{\kappa^3 \ln 2}\frac{m_{3/2}}{v} \quad (\leq 1)\,. \tag{27}$$

Thus we obtain the ellipticity parameter of inflaton trajectory after inflation as

$$\epsilon \simeq \frac{n_{\rm inf}}{m_{\rm inf} v^2} \simeq \frac{H_{\rm inf}}{m_{\rm inf}}\xi, \tag{28}$$

where the factor of $H_{\rm inf}/m_{\rm inf}$ comes from the difference of the time scale of inflaton dynamics between the eras during and after inflation.

When we consider the gravitino mass of order 100 TeV, the spectral index can be consistent with the observed value for the case of $\lambda \simeq 3 \times 10^{-3}$ and $v \simeq 4 \times 10^{15}$ GeV for the final phase value of order unity [25]. This implies that the mass of inflaton and the Hubble parameter just after inflation are given by

$$m_{\rm inf} = \lambda v \simeq 1.2 \times 10^{13}\ {\rm GeV} \tag{29}$$

$$H_{\rm inf} \simeq \frac{\lambda v^2}{2\sqrt{3}M_{\rm Pl}} \simeq 5.7 \times 10^{9}\ {\rm GeV}. \tag{30}$$

From Eq. (20), we find that the reheating temperature needs to be as high as

$$6 \times 10^{13}\ {\rm GeV}\epsilon^{-2/7}\left(\tilde{G}_F M_{\rm Pl}^2\right)^{-2/7}, \tag{31}$$

to explain the observed baryon asymmetry. Note that for the above parameters the ellipticity parameter ϵ is of order 10^{-4}, so that we need $\tilde{G}_F M_{\rm Pl}^2 \simeq 10^4$ (i.e., $c \simeq 10^4$) to explain the observed baryon asymmetry. In addition, the reheating temperature cannot be as high as 10^{13} GeV in the hybrid inflation model. This is because couplings between inflaton and other particles have to be suppressed by κ ($\ll 1$) in order not to affect the inflaton potential and the nonperturbative enhancement of decay process called preheating is suppressed after inflation in the case of the rotating inflaton [26]. Thus, we conclude that both the ellipticity parameter ϵ and the reheating temperature are too small to account for the observed baryon abundance in the hybrid inflation model.

B. Chaotic inflation

As another example, let us focus on a chaotic inflation model proposed in Ref. [27], where the Kähler potential respects a shift symmetry of the inflaton field,

$$\phi \to \phi + i\alpha, \tag{32}$$

where α is the transformation parameter. The Kähler potential is given by

$$\begin{aligned} K = c'(\phi+\phi^*) + \frac{1}{2}(\phi+\phi^*)^2 + |X|^2 \\ + |\chi|^2 + \frac{c}{2}(\phi+\phi^*)^2|\chi|^2, \end{aligned} \tag{33}$$

where X is a stabilizer field. The relevant interactions of the Lagrangian are then given by

$$\begin{aligned} \mathcal{L}_{\rm int} =& \frac{c'}{2}\partial_\mu {\rm Im}[\phi]\bar{\chi}\bar{\sigma}^\mu\chi \\ & -2{\rm Re}[\phi]\partial_\mu{\rm Im}[\phi]\left[2c{\rm Im}\left(\tilde{\chi}\partial^\mu\tilde{\chi}^*\right) + \left(c - \frac{1}{2}\right)\bar{\chi}\bar{\sigma}^\mu\chi\right], \end{aligned} \tag{34}$$

where we have rescaled the fields to obtain the canonical kinetic terms. The shift symmetry is assumed to be explicitly broken in the superpotential,

$$W = m_{\rm inf}\phi X, \tag{35}$$

where the R-charge assignment is $R[X] = 2$ and $R[\phi] = 0$. When the constant c' is nonzero, the real part of ϕ has a VEV of order c' during inflation. Thus the inflaton starts to rotate in the complex plane after inflation such as

$${\rm Re}[\phi] \approx c'|\phi|\sin m_{\rm inf}t \tag{36}$$

$${\rm Im}[\phi] \approx |\phi|\cos m_{\rm inf}t. \tag{37}$$

Note here that, in addition to the $U(1)_R$ symmetry, the scalar potential has another (approximate) global U(1) symmetry for which ϕ and X have the same magnitude charge but opposite sign. This implies that the inflaton number is induced after inflation:

$${\rm Re}[\phi]\partial_0{\rm Im}[\phi] \simeq c' m_{\rm inf}|\phi|^2 \simeq c'\frac{H^2}{m_{\rm inf}}. \tag{38}$$

That is, the ellipticity parameter is given by $\epsilon \simeq c'$, which is expected to be of order unity, and the effective coupling is given by $(\tilde{G}_F M_{\rm Pl}^2) \approx c$.

The mass of inflaton is determined by the COBE normalisation such as $m_{\rm inf} \simeq 10^{13}$ GeV, which implies that the Hubble parameter just after inflation is given by 6×10^{12} GeV. Thus the observed baryon asymmetry can be explained when the reheating temperature is as large as

$$T_{\rm RH} \simeq 3 \times 10^{13}\ {\rm GeV}\left(\tilde{G}_F M_{\rm Pl}^2\right)^{-2/7}, \tag{39}$$

where we assume $\epsilon = 1$. Note that $(\tilde{G}_F M_{\rm Pl}^2)$ ($\approx c$) is expected to be of order unity. To obtain such high reheating temperature, we introduce a superpotential of [27]

$$W_{\rm RH} = yXH_uH_d, \tag{40}$$

where H_u and H_d are a pair of Higgs doublets in the MSSM sector. This leads to the coupling between the inflaton and the Higgs doublets as $\mathcal{L}_{\rm RH} \sim y m_{\rm inf} \phi H_u H_d$, leading to the reheating temperature given by[2]

$$T_{\rm RH} \sim 10^{13}\ {\rm GeV} \left(\frac{y}{0.1}\right)\left(\frac{m_{\rm inf}}{10^{13}\ {\rm GeV}}\right)^{1/2}. \qquad (41)$$

Therefore, in the chaotic inflation model our scenario of spontaneous baryogenesis works naturally and can explain the observed baryon abundance without any fine-tunings of the parameters.

IV. DISCUSSION AND CONCLUSIONS

In this paper we have proposed a scenario of spontaneous baryogenesis from asymmetric inflaton via current-current interactions between the inflaton and $B-L$ numbers. Such interactions are naturally present in supergravity theories. The CP asymmetric part of the inflaton number induces an effective chemical potential of $B-L$ number, which biases the $B-L$ asymmetry in the equilibrium state. If the $B-L$ number is broken in the plasma, a non-zero $B-L$ asymmetry is induced. We have shown that the observed baryon abundance can be explained by this mechanism when the reheating temperature is as large as 10^{13} GeV. We have also studied concrete inflation models in supergravity to see if the above mechanism can be implemented, and we have found that our mechanism naturally explain the observed baryon abundance without fine-tunings in the chaotic inflation model.

The quadratic chaotic inflation predicts a rather large tensor-to-scalar ratio r, which is now strongly disfavored by the CMB observations [28, 29]. The tensor-to-scalar ratio can be easily reduced by introducing higher order terms of the inflaton, in which case the inflaton potential is given by a polynomial function [30, 31]. Our spontaneous baryogenesis scenario works similarly in this case.

It is known that the spontaneous baryogenesis in the slow-roll regime generically leads to sizable isocurvature perturbations [13]. This makes the spontaneous baryogenesis incompatible with high-scale inflation. In the case of chaotic inflation studied in the previous section, both the real component of ϕ as well as the stabilizer field can have a mass of order the Hubble parameter during inflation, and so there is no light degrees of freedom during inflation other than the inflaton. Thus, no isocurvature perturbations are induced.

In general, high reheating temperature $T_{RH} \sim 10^{13}$ GeV is required for successful baryogenesis in our scenario, and gravitinos are copiously produced from thermal scattering. On the other hand, non-thermal gravitino production can be suppressed at such high reheating temperature [32–34]. If the gravitino is lighter than of order 100 TeV, its decay destroys light elements, altering their abundances in contradiction with observations. If it is heavier than 100 TeV, it decays before the BBN epoch but the LSPs may be overproduced. In order to avoid the overproduction of the LSPs, we may assume that the R-parity is violated and the LSP decays before the BBN epoch. In this case, we require another dark matter candidate such as axion.

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