Affleck-Dine baryogenesis with modulated reheating

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Modulated reheating scenario is one of the most attractive models that predict possible detections of not only the primordial non-Gaussianity but also the tensor fluctuation through future CMB observations such as the Planck satellite, the PolarBear and the LiteBIRD satellite experiments. We study the baryonic-isocurvature fluctuations in the Affleck-Dine baryogenesis with the modulated reheating scenario. We show that the Affleck-Dine baryogenesis can be consistent with the modulated reheating scenario with respect to the current observational constraint on the baryonic-isocurvature fluctuations.

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I. INTRODUCTION

The observed cosmological perturbation is generated from a primordial curvature perturbation. It is known that the curvature perturbation can be time-independent after the cosmic time approximately becomes one second ($t \gtrsim 1$ sec) because the cosmic fluid had been perfectly radiation-dominated until that time [1]. We expect the corresponding perturbations can be generated from the vacuum fluctuation of a light field, which requires the mass of the field to be lighter than the Hubble parameter $H$ during inflation [2]. In original single-field slow-roll inflation paradigm, the curvature perturbation mainly originated from the perturbation of the so-called inflaton field which induces the primordial inflation. Then the curvature perturbation had been already produced at the initial epoch during inflation and is constant. In this case, a deviation from the Gaussian statistics of the fluctuation (so called non-Gaussianity) is small since the perturbation does not change the relation between the energy density $\rho$ and the pressure density $P$ very much through the cosmic history, and the non-linearity parameter to express the degree of the non-Gaussianity should be the order of the slow-roll parameters $f_{NL} \sim \epsilon$ and/or $\eta \sim \mathcal{O}(10^{-2})$. Then it means that we will not be able to detect the non-Gaussianity.

When we consider other paradigms by assuming other light field $\sigma$, the curvature perturbation may originate mainly from the perturbation of $\sigma$. Although the curvature perturbation from $\delta \sigma$ might be initially negligible, it will grow later even after inflation ends because it generates a non-adiabatic pressure perturbation. This mechanism works at any epochs except for complete radiation domination or complete matter domination. In this case we should be able to detect those different types of non-Gaussianity in future.

So far a variety of models which produce a non-adiabatic pressure perturbation and generate non-Gaussianity even in terms of slow-roll inflation have been reported in this direction; during multi-field inflation [3, 4], before a second reheating through the curvaton mechanism [5], at the end of inflation [6-10], during modulated reheating [11-15] or modulated preheating [16, 17], at a modulated phase transition [18, 19], and a modulated trapping [20]. By observing the non-Gaussianity, we will be able to discern a better model from other ones. In this paper, especially we focus on modulated reheating because of some attractive reasons mentioned later.

The latest constraint on the non-linearity parameter $f_{NL}$ from the WMAP 7-year data is $f_{NL} = 32 \pm 21$ (68\%C.L.) [21] for the local type of non-Gaussianity, which means a null detection of the non-Gaussianity. On the other hand, the Planck satellite is expected to reach $\Delta f_{NL} \sim 5$. [22] Therefore we may be able to detect the non-Gaussianity in near future within 5 years.

Another good indicators to discern a model from others must be the tensor to scalar ratio $r$ for the perturbation. So far WMAP 7-year has reported only its upper bound $r \lesssim 0.2$ [21], which does not exclude even the standard single-field quadratic chaotic inflation model. It is expected that the PLANCK satellite will
be able to observe $r \sim 0.1 \ [22, 23]$. In addition, it should be quite exciting that future ground-based detectors QUIET+PolarBear will reach $r \sim \mathcal{O}(0.01) \ [23]$, and KEK’s future CMB satellite experiment, LiteBIRD will observe $r \sim \mathcal{O}(10^{-3}) \ [23, 26]$. Considering these current and future situations for the detectability of non-Gaussianity and the tensor fluctuation, for comparison in future it should be necessary to keep those known models in proper order with respect to the predicted values of $r$ and $f_{NL}$, respectively. However, unfortunately it might be known that there exist few models which predict both the large non-Gaussianity of $f_{NL} \sim \mathcal{O}(10)$ and the large tensor to scalar ratio of $r \sim \mathcal{O}(0.1)$. Therefore, it should be required for theorists in advance to build models which predict both the large $f_{NL}$ and large $r$ enough to be detectable in near future.

For this purpose, the modulated reheating model must be attractive because it possibly produces large non-Gaussianity and can predict a large tensor to scalar ratio if the inflation scale is sufficiently large like chaotic inflation with its Hubble parameter $H \sim 10^{13}$ GeV.

On the other hand, we also have to check consistencies of the modulated reheating scenario with other observational constraints. As will be discussed later in detail, in the modulated reheating scenario the curvature perturbation is effectively governed by the fluctuation of the reheating temperature after inflation $\delta T_R/T_R$. Since most class of viable baryogenesis scenarios in modern cosmology depend on the reheating temperature, the modulated reheating may induce a large baryonic-isocurvature fluctuation.\footnote{Some class of scenarios for dark-matter production also depend on the reheating temperature such as gravitino thermal/non-thermal production. In this case the modulated reheating is severely constrained by observations of the cold dark matter (CDM)-isocurvature fluctuation \[25, 26\]. In this paper we are assuming a dark matter such as the lightest neutralino which was decoupled from the thermal bath with its appropriate thermal relic density to fit the observation. Then the modulated reheating does not produce a sizable CDM-isocurvature fluctuation.}

In this paper we consider baryogenesis in models with supersymmetric (SUSY) extension of standard model, especially so-called Affleck-Dine (AD) mechanism \[29, 30\], which is naturally realized even in the Minimal Supersymmetric Standard Model (MSSM)\footnote{See also Ref. \[12\] for another mechanism of baryogenesis in the SUSY cosmologies and its compatibility with the modulated scenario.} and agrees with observations in broad parameter regions. Since good candidates for the light scalar field $\sigma$ could be found in SUSY \[11\] or supergravity (the local theory of SUSY), this direction of discussion should be naturally motivated. As we have already raised the question, however, it might be nontrivial if the AD baryogenesis is consistent with the modulated reheating scenarios because the produced baryon number sometimes depends on the reheating temperature in the normal parameter space. Thus we have a strong motivation to search the allowed parameter region for the AD baryogenesis to avoid the constraint on the baryonic-isocurvature fluctuation in the modulated reheating scenario.

This paper is organized as follows. In Sec. II we show the basic picture of the modulated reheating scenarios and the conditions for the decay rate of the inflaton field where the large non-Gaussianity can be predicted. The AD barygenesis is outlined in Sec. III. In Sec. IV we look for the parameter space to agree with the observations and also discuss a possible isocurvature fluctuation that the AD barygenesis may originally have. Sec. V is devoted to conclusions.

\section{II. MODULATED REHEATING SCENARIO}

Here, we give a brief review of the modulated reheating scenario \[11\]. In such scenario, we consider the decay rate of the inflaton, $\Gamma$, depending on a light scalar field, $\sigma$, which has a quantum fluctuation during inflation, that is, $\Gamma = \Gamma(\sigma)$. The $e$-folding number $N = \int d\ln a$, where $a$ is a scale factor, measured between the end of inflation at $t = t_{inf}$ and a time after the end of the complete reheating, $t_c$, is given by

$$N = \ln \left( \frac{a(t_c)}{a(t_{inf})} \right)$$

$$= \ln \left( \frac{a(t_{reh})}{a(t_{inf})} \right) + \ln \left( \frac{a(t_c)}{a(t_{reh})} \right), \tag{1}$$

where $t_{reh}$ represents a time at $d\ln a/dt = H = \Gamma$.

Let us consider the quadratic inflaton potential, $V(\phi) \propto \phi^2$. In such case, during the inflaton oscillating phase after the inflation, the energy density of the Universe relying on the inflaton decays as $\rho \propto a^{-3}$ and the Hubble parameter, $H$, evolves as $H \propto \rho^{1/2}$. Since after the complete reheating the energy density of the Universe is dominated by the radiation ($\rho \propto a^{-4}$ and $H \propto a^{-2}$), the $e$-folding number given by Eq. \[1\] is rewritten as

$$N = \ln \left( \frac{a(t_{reh})}{a(t_{inf})} \right) + \ln \left( \frac{a(t_c)}{a(t_{reh})} \right)$$

$$= -\frac{1}{6} \ln \left( \frac{\Gamma}{H(t_{inf})} \right) + \frac{1}{2} \ln \left( \frac{H(t_{inf})}{H(t_c)} \right), \tag{2}$$

where we have used $H(t_{reh}) = \Gamma$. The fluctuation of $\sigma$ induces the modulated reheating and hence the fluctuation of the $e$-folding number is given by

$$\delta N = -\frac{1}{6} \delta \Gamma(\sigma) + \frac{1}{2} \Gamma' \delta \sigma = -\frac{1}{6} \Gamma' \delta \sigma, \tag{3}$$

where $\Gamma'(\sigma) = d\Gamma(\sigma)/d\sigma$. In terms of the reheating temperature $T_R \propto \Gamma^{1/2}$, the above expression can be rewritten as

$$\delta N = -\frac{1}{3} \frac{\delta T_R}{T_R}. \tag{4}$$
Based on $\delta N$ formalism $^{31}$-$^{37}$, the curvature perturbation on the uniform energy density hypersurface, $\zeta$, on super-horizon scales is given by the perturbation of the $e$-folding number as $\zeta \approx \delta N$ and hence we find that in the modulated reheating scenario the curvature perturbation can be generated due to the fluctuation of the decay rate of the inflaton. Up to the second order, we have

$$\zeta \approx \delta N = -\frac{1}{6} \frac{\Gamma'}{T} \left[ \delta \sigma + \frac{1}{2} \left( \frac{\Gamma''}{\Gamma} - \frac{\Gamma'}{T} \right) \delta \sigma^2 \right].$$

Hence the power spectrum and the non-linearity parameter $f_{NL}$ defined as

$$\langle \zeta(k)\zeta(k') \rangle \equiv (2\pi)^3 \frac{2\pi^2}{k^3} P(k)\delta^{(3)}(k+k'),$$

$$\zeta = \zeta_{\text{lin}} + \frac{3}{5} f_{NL} \zeta_{\text{lin}},$$

are respectively given by

$$P(k) = \left( \frac{1}{6} \frac{\Gamma'}{T} \right)^2 \left( \frac{H_{\text{inf}}}{2\pi} \right)^2,$$

$$f_{NL} = 5 \left( 1 - \frac{\Gamma''}{\Gamma T} \right).$$

As an example, let us consider an interaction term in Lagrangian as $L \supset -g(\sigma)\bar{\psi}\psi\bar{\phi}$ where $g(\sigma)$ is a coupling constant and $\psi$ is a light fermion which interacts with the standard particles and its energy finally goes to thermal bath. Then we obtain the decay rate of the inflaton $\Gamma$ as

$$\Gamma(\sigma) \simeq \frac{g(\sigma)^2}{8\pi} m_\phi,$$

where $m_\phi$ is the mass of the inflaton. Assuming a $\sigma$-dependence of $g$ as

$$g(\sigma) = g \left[ 1 + \lambda \left( \frac{\sigma}{M_{\text{cut}}} \right)^2 \right],$$

where $\lambda(\sim O(1))$ is some constant parameter and $M_{\text{cut}}$ is a cutoff scale of the effective theory $^{13}$-$^{14}$-$^{16}$-$^{19}$, the non-linearity parameter $f_{NL}$ is given by

$$f_{NL} \simeq -10^2 \lambda^{-1} \left( \frac{\sigma}{M_{\text{cut}}} \right) \left( \frac{0.1}{0.1} \right)^{-2},$$

where we have assumed that $\sigma/M_{\text{cut}} \ll 1$. The power spectrum is also obtained as

$$P(k) \simeq 10^{-10} \times$$

$$\left( \frac{H_{\text{inf}}}{10^{13}\text{GeV}} \right)^2 \left( \frac{M_{\text{cut}}}{10^{16}\text{GeV}} \right)^{-2} \lambda^2 \left( \frac{\sigma}{M_{\text{cut}}} \right) \left( \frac{0.1}{0.1} \right)^2,$$

which can be marginally comparable to the observational value $P_{\text{obs}}(k) \approx 3 \times 10^{-10}$ $^{21}$ or subdominant for normal scales of $M_{\text{cut}}$ ($\sim 10^{16}$ GeV $- 10^{18}$ GeV)$^3$.

Hence, in the modulated reheating scenario we can easily find that the large non-Gaussianity ($f_{NL} \sim O(10)$) can be generated even for the high energy inflation with the large $r = 16\epsilon \sim 0.1(H_{\text{inf}}/10^{13}\text{GeV})^2$ when the model parameters are appropriately chosen.

III. AFFLECK-DINE MECHANISM

AD mechanism $^{23}$-$^{30}$ has been known as one of the powerful candidates for the successful baryogenesis mechanism. It can be realized by taking advantage of a flat direction along which scalar potential vanishes in the global SUSY limit. Hereafter we call the complex scalar field that parameterizes the flat direction as $\Phi$ field and assume that it carries non-zero baryon charge $\beta$.

Though the scalar potential for the $\Phi$ field vanishes in the global SUSY limit, it is lifted by non-renormalizable terms, the SUSY-breaking effect and some other effects. Let us consider a non-renormalizable superpotential for the $\Phi$ field given by

$$W_{\text{nr}} = \frac{\Phi^{n+3}}{(n+3)M_*^n},$$

where $M_*$ is the cut-off scale and the positive integer $n$ depends on the flat direction. Including the SUSY breaking effect, the induced scalar potential reads

$$V = V_\Phi + \frac{\Phi^{2n+4}}{M_*^{2n}} + \left( \frac{a_B m_{3/2}}{M_*} \Phi^{n+3} + \text{h.c.} \right),$$

where $m_{3/2}$ is a gravitino mass and $a_B$ is a complex numerical factor whose amplitude is of order of unity. $V_\Phi$ is the soft SUSY breaking effect that depends on the SUSY breaking mechanism. The second term is the $F$-term that comes from non-renormalizable operator $W_{\text{nr}}$. The last term represents the interaction between non-renormalizable operator and the SUSY breaking sector coming from supergravity effect, which breaks the $U(1)$ baryon symmetry and is called as the $A$-term.

During and after inflation, the AD field acquires the Hubble induced mass from the interaction between the AD field and the inflaton through the supergravity effect, which can be negative,

$$V_H = -c_H H^2 |\Phi|^2,$$
where $c_H$ is a positive numerical factor of order of unity.

We must note that even before reheating there can be thermal plasma just after the end of inflation as a subdominant component of the universe. Its temperature can be expressed as [38]

$$T \simeq (HM_G T_R^4)^{1/4},$$

(17)

where $M_G$ is the reduced Planck scale. Since the AD field can have interaction with the thermal plasma directly or indirectly, it acquires a thermal potential [39],

$$V_{\text{thermal}} \sim \begin{cases} h^2 T^2 |\Phi|^2 & (h|\Phi| \ll T), \\ \alpha_g^2 T^4 \log \left( \frac{h^2 |\Phi|^2}{T^2} \right) & (h|\Phi| \gg T), \end{cases}$$

(18)

in addition to above terms. Here $h$ is the Yukawa or the gauge coupling constant of the AD field, and $\alpha_g \equiv g^2 / 4 \pi$ represents the gauge coupling constant. The upper term is the thermal mass from the thermal plasma. On the other hand, the lower one (called the thermal logarithmic term) represents the two-loop finite temperature effects coming from the running of the gauge coupling with the non-zero value of the AD field. 5

Thus, the AD field evolves with the effective potential,

$$V_{\text{eff}} = V + V_H + V_{\text{thermal}}.$$  

(19)

During and after inflation, when the Hubble parameter $H$ is sufficiently large, the AD field settles down to the time-dependent potential minimum,

$$|\Phi| \simeq (HM_s^n)^{1/(n+1)},$$

(20)

and traces its evolution. Note that there can be several non-renormalizable operators for the AD field but only the one with the smallest $n$ determines the dynamics of the AD field. Thus hereafter we consider only smaller $n$ ($n \leq 3$).

Let us consider the evolution of the AD field further. As the Hubble parameter decreases, the Hubble induced mass also gets small. Then, when $H_{\text{osc}}^2 \simeq |V_{\text{eff}}'|$, the AD field (more precisely its radial component) starts to oscillate around the origin. Here the dash denotes the derivative with respect to $\phi \equiv \sqrt{2}|\Phi|$, and hereafter the subscript “osc” indicates that the parameter or the variable is evaluated at the onset of the AD field oscillation.

At the onset of the oscillation, the AD field acquires an angular momentum due to the $A$-term, which represents the baryon asymmetry of the Universe $n_B$,

$$n_B(t_{\text{osc}}) \simeq \beta m_{3/2} (H_{\text{osc}} M_s^n)^{2/(n+1)} \sin(n\theta_{\text{inf}} + \alpha),$$

(21)

where $\theta_{\text{inf}}$ and $\alpha$ are the phases of $\Phi$ during inflation and the constant $a_B$ in the third term of Eq. (15), respectively. Just after the onset of the AD field oscillation, $\alpha^3 n_B$ is conserved since the CP-violating $A$-term comes to ineffective quickly. This is because the AD field value continues decreasing with time during the field oscillation due to the cosmic expansion. Since the entropy density decreases as $s \propto a^{-3}$ after the reheating if there is no late-time entropy production, the baryon-to-entropy ratio $n_B/s$ is conserved. Thus its present value is estimated as

$$\left(\frac{n_B}{s}\right)_0 \simeq \frac{\beta m_{3/2} T_R}{M_G^2 H_{\text{osc}}} (H_{\text{osc}} M_s^n)^{2/(n+1)} \sin(n\theta_{\text{inf}} + \alpha).$$

(22)

Note that there are four scenarios according to which term in the potential drives the AD field oscillation. Thus $H_{\text{osc}}$ is different from each other [40]. The first scenario is that the AD field oscillation is driven by the zero-temperature potential $V_S$. In this case, the Hubble parameter at the onset of the AD field oscillation is

$$H_{\text{osc}} \simeq m_0(|\Phi|_{\text{osc}}),$$

(23)

where $m_0(|\Phi|) \equiv V_S''(|\Phi|)$. The second one is that it is driven by the thermal logarithmic term and

$$H_{\text{osc}} \simeq \left( \frac{M_G^2 T_R^4}{M_s^{2n}} \right)^{1/(n+3)}.$$

(24)

The third possibility is that it is driven by the thermal mass and

$$H_{\text{osc}} \simeq (h^4 M_G T_R^2)^{1/3}.$$ 

(25)

When the Yukawa coupling $h$ is rather small, there is a discrepancy in the thermal potential at $|\Phi| \simeq T/h$. Thus, if $n < 3$, it is possible that AD field start to oscillate immediately after the AD field value becomes $|\Phi| \simeq T/h$. In this case, the Hubble parameter at the onset of the AD field oscillation,

$$H_{\text{osc}} \simeq \left( \frac{M_G^{n+1} T_R^{2n+1}}{h^{2n+1} M_s^{2n}} \right)^{1/(3-n)}.$$ \hspace{1cm} (26)

This classification is essential to estimate the baryonic-isocurvature perturbation as we will see in the next section.

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4 In the case where $c_H$ is negative, the AD field falls into the origin quickly and does not affect cosmology.

5 The signature of the thermal logarithmic term can be both positive or negative depending on the flat direction. For example, the $LH_\Phi$ flat direction receives a positive thermal logarithmic term [40]. Here we consider only positive thermal logarithmic terms. The case with negative thermal logarithmic term is discussed in Ref. [11] in a different context.
IV. BARYONIC-ISOCURVATURE FLUCTUATION

Let us consider the baryonic-isocurvature fluctuation $S_B$, which is commonly defined as

$$S_B \equiv \frac{\delta n_B}{n_B} - \frac{\delta s}{s} = \frac{\delta (n_B/s)}{n_B/s}. \quad (27)$$

In the case where the baryon-to-entropy ratio depends on the reheating temperature, the baryonic-isocurvature fluctuation can be also generated in the modulated reheating scenario [27]. If we assume $n_B/s \propto T_R^2$, we have

$$S_B = \frac{\delta (n_B/s)}{n_B/s} = p\frac{\delta T_R}{T_R}. \quad (28)$$

Hence, in the case where the curvature perturbation originates mainly from the modulated reheating mechanism, from Eq. (4) and Eq. (28), we have

$$S_B = -3p\zeta. \quad (29)$$

The current observational limit for the anti-correlated baryonic-isocurvature fluctuation is roughly given by $|S_B/\zeta| \lesssim O(0.1)$ [21, 42] and hence it means that the models with $p \gtrsim O(1)$ are conflict with current observations. From Eq. (22), we find that in general the present baryon-to-entropy ratio generated by the AD mechanism strongly depends on the reheating temperature. Thus it seems that the AD mechanism and the modulated reheating scenario are incompatible.

However, when $n = 1$ and $H_{\text{osc}} \simeq (M_G T_R^2/M_*)^{1/2}$ (Eq. 24) or $n = 3$ and $H_{\text{osc}} \simeq (h^4 M_G T_R^2)^{1/3} (h \simeq 10^{-5})$ (Eq. 25), the dependence of the present baryon-to-entropy ratio on the reheating temperature is canceled. If such situations are realized, baryonic-isocurvature fluctuation is not generated from the modulated reheating. It is non-trivial whether there is a parameter space in which the Hubble parameter at the onset of the AD field oscillation comes to the value above and the baryon-to-entropy ratio can be explained the present value, $(n_B/s)_0 \simeq O(10^{-10})$. In fact, there is two sets of parameter spaces in the gravity mediated SUSY-breaking mechanism 10: (i) $n = 1, M_* \simeq 10^{22-23}$ GeV, and $T_R \gtrsim 10^7$ GeV, and (ii) $n = 3, M_* \simeq 10^{16}$ GeV, and $T_R \simeq 10^6$ GeV. In the case (i) this cut-off scale $M_*$ suggests that small lightest neutrino mass for the LH$_3$ flat direction 10 and in the case (ii) it coincides with the GUT scale. In these cases, the AD mechanism and the modulated reheating scenario are compatible.

Here we comment on another source of the baryonic-isocurvature perturbation. In the absence of the Hubble induced $A$-term, the baryonic-isocurvature perturbation is generated from the fluctuation for the angular component of the AD field [44]. The magnitude of the fluctuation is given by

$$\delta \theta \simeq \frac{H_{\text{inf}}}{2\pi |\Phi_{\text{inf}}|} \simeq \frac{1}{2\pi} \left(\frac{H_{\text{inf}}}{M_*}\right)^{n/(n+1)}, \quad (30)$$

where $H_{\text{inf}}$ and $|\Phi_{\text{inf}}|$ are the Hubble parameter and the AD field value during inflation. Thus, from Eq. (22) and Eq. (27), we have the net baryonic-isocurvature perturbation as

$$S_B \simeq \frac{n}{2\pi \cot(n\theta_{\text{inf}} + \alpha)} \left(\frac{H_{\text{inf}}}{M_*}\right)^{n/(n+1)} + p\frac{\delta T_R}{T_R}. \quad (31)$$

By considering the observational constraint on the uncorrelated baryonic isocurvature mode ($|S_B/\zeta| \lesssim O(1)$ [21, 42]) originated from the fluctuation for the angular component of the AD field (e.g., see Ref. [43]) we find that the Hubble parameter during inflation should be restricted to be (i) $H_{\text{inf}} \lesssim 10^{13}$ GeV and (ii) $H_{\text{inf}} \lesssim 10^9$ GeV, respectively. Thus in the case (i) the chaotic inflation with $H_{\text{inf}} \sim 10^{13}$ GeV is still allowed by the observation in terms of the baryonic-isocurvature fluctuation. In such case, we can expect the large tensor-to-scalar ratio at the detectable level in the future experiments, for examples, such as Planck satellite, PolarBear or LiteBIRD satellite.

V. CONCLUSION

Recently, the primordial non-Gaussianity and the tensor fluctuations have been focus of attention to provide the information about the physics of the early Universe, especially, the origin of the observed cosmological perturbations. There would be a lot of the scenarios generating large non-Gaussianity, and the ones generating large tensor-to-scalar ratio, respectively at the detectable level in the future experiments independently of each other.

On the other hand, however, the modulated reheating scenario is quite attractive since it can produce the large non-Gaussianity, and would simultaneously predict the large tensor to scalar ratio when the inflation scale were large like the case of chaotic inflation.

Of course, we need to check the consistencies of the modulated reheating scenario as a mechanism of generating primordial curvature fluctuations with other phenomena in the early Universe. In this paper, we focused on the baryon asymmetry in the Universe. In particular, we consider the one of the most promising candidates for baryogenesis, the AD mechanism.

We have shown that the modulated reheating scenario is consistent with AD baryogenesis in some sets of model

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6 In these reheating temperatures, the gravitino problem can be avoidable [13].

7 Here we assume $\cot(n\theta_{\text{inf}} + \alpha) \simeq O(1) - O(0.1)$. This is valid when there is no fine-tuning for the initial phase of the AD field.
parameters, which can be motivated from the physics at the high-energy scales, even if the observational constraint for the isocurvature fluctuation and the gravitino problem is imposed on the prediction. This conclusion is not changed even if we consider the uncorrelated baryonic-isocurvature mode originated from the fluctuation of the angular component in the AD field. Therefore we will be able to discern this model from others in principle when we detect a large $f_{NL}$, and a large $r$ in near future.

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