Iso-curvature fluctuations in modulated reheating scenario

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Abstract

Modulated reheating scenario is one of the most attractive models that predict possible detections of primordial non-Gaussianity through future CMB observations such as the Planck satellite. We study the baryonic-isocurvature fluctuations in the Affleck-Dine baryogenesis with the modulated reheating scenario. We show that the simple Affleck-Dine baryogenesis would be incompatible with the modulated reheating scenario with respect to the current observational constraint on the baryonic-isocurvature fluctuations, like a gravitino dark matter scenario.

1 Introduction

Primordial non-Gaussianity of curvature fluctuation fluctuations is one of well discussed topics recently. From the recent result from WMAP 7-year data, the so-called local type non-linearity parameter $f_{\rm NL}$ is constrained as $-10 < f_{\rm NL} < 74$ [1]. If future observations confirm such a large value of $f_{\rm NL}$, we need some mechanism which generates large non-Gaussian primordial fluctuations. The modulated reheating scenario [2] , where a scalar field other than the inflaton is responsible for primordial fluctuations through the inhomogeneous reheating, is interesting mechanisms generating large non-Gaussianity.

Of course, it is important to check consistencies of the modulated reheating scenario with other observational constraints. 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. Some class of scenarios for dark-matter production also depend on the reheating temperature such as gravitino thermal/non-thermal production. In Ref. [3, 4], the authors shown that in this case the modulated reheating is severely constrained by observations of the cold dark matter (CDM)-isocurvature fluctuation.

Here, we consider baryogenesis in models with supersymmetric (SUSY) extension of standard model, especially so-called Affleck-Dine (AD) mechanism [5, 6], which is naturally realized even in the Minimal Supersymmetric Standard Model (MSSM) and agrees with observations in broad parameter regions [7]. Since good candidates for the light scalar field σ could be found in SUSY or supergravity (the local theory of SUSY), this direction of discussion should be naturally motivated.

2 Modulated Reheating Scenario

Here, we give a brief review of the modulated reheating scenario. In such scenario, we consider the decay rate of the inflaton, Γ, depending on a light scalar field, σ, which has a quantum fluctuation during inflation, that is, $\Gamma = \Gamma(\sigma)$.

In order to evaluate the curvature perturbation generated in the modulated reheating scenario, let us consider the e-folding number $N = \int d\ln a$, where a is a scale factor, measured between the end of inflation at $t = t_{\rm inf}$ and a time after the end of the complete reheating, $t_{\rm c}$. This can be written as

$$\begin{aligned} N &= \ln\left(\frac{a(t_{\rm c})}{a(t_{\rm inf})}\right) \\ &= \ln\left(\frac{a(t_{\rm reh})}{a(t_{\rm inf})}\right) + \ln\left(\frac{a(t_{\rm c})}{a(t_{\rm reh})}\right) \ , \end{aligned} \tag{1}$$

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where $t_{\rm reh}$ represents a time at $d\ln a/dt = H = \Gamma$. For the quadratic inflaton potential, $V(\phi) \propto \phi^2$, 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

$$N = \ln\left(\frac{a(t_{\rm reh})}{a(t_{\rm inf})}\right) + \ln\left(\frac{a(t_{\rm c})}{a(t_{\rm reh})}\right) = -\frac{1}{6}\ln\left(\frac{\Gamma}{H(t_{\rm inf})}\right) + \frac{1}{2}\ln\left(\frac{H(t_{\rm inf})}{H(t_{\rm c})}\right) , \tag{2}$$

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

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

where $\Gamma'(\sigma) \equiv d\Gamma(\sigma)/d\sigma$. Based on δN formula, taking the final hypersurface at $t = t_{\rm c}$ to be a uniform energy density one, we have

$$\zeta(t_{\rm c}) = \delta N(t_{\rm c}, t_{\rm ini}) = -\frac{1}{6}\frac{\Gamma'}{\Gamma}\delta\sigma(t_{\rm ini}) , \tag{4}$$

where $\delta\sigma(t_{\rm ini})$ is the fluctuation of the field σ on the flat hypersurface and we have assumed that $\delta\sigma$ is almost constant and also almost Gaussian fluctuation after the horizon crossing time. In terms of the reheating temperature $T_R \propto \Gamma^{1/2}$, the above expression can be rewritten as

$$\delta N = -\frac{1}{3}\frac{\delta T_R}{T_R} . \tag{5}$$

Up to the second order, we have

$$\zeta \approx \delta N = -\frac{1}{6}\frac{\Gamma'}{\Gamma}\left[\delta\sigma + \frac{1}{2}\left(\frac{\Gamma''}{\Gamma'} - \frac{\Gamma'}{\Gamma}\right)\delta\sigma^2\right] . \tag{6}$$

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

$$\langle\zeta(\mathbf{k})\zeta(\mathbf{k}')\rangle \equiv (2\pi)^3\frac{2\pi^2}{k^3}\mathcal{P}(k)\delta^{(3)}(\mathbf{k}+\mathbf{k}') , \tag{7}$$

$$\zeta = \zeta_{\rm lin} + \frac{3}{5}f_{\rm NL}\zeta_{\rm lin}^2 , \tag{8}$$

are respectively given by

$$\mathcal{P}(k) = \left(\frac{1}{6}\frac{\Gamma'}{\Gamma}\right)^2\left(\frac{H_{\rm inf}}{2\pi}\right)^2 , \tag{9}$$

$$f_{\rm NL} = 5\left(1 - \frac{\Gamma\Gamma''}{\Gamma'^2}\right) . \tag{10}$$

Hence, in the modulated reheating scenario we can easily get the large non-Gaussianity ($f_{\rm NL} \sim \mathcal{O}(10)$) by considering appropriate form of $\Gamma(\sigma)$ and also the power spectrum of primordial curvature fluctuations which is consistent with the current cosmological observations.

3 Affleck-Dine Baryogenesis Scenario

AD mechanism [5, 6] 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 AD field Φ and assume that it carries non-zero baryon charge β.

Though the scalar potential for the AD 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 AD field given by

$$W_{\rm nr} = \frac{\Phi^{n+3}}{(n+3)M_*^n}, \tag{11}$$

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

$$V = V_{\not{S}} + \frac{|\Phi|^{2n+4}}{M_*^{2n}} + \left(\frac{a_B m_{3/2}}{M_*^n}\Phi^{n+3} + {\rm h.c.}\right)\,, \tag{12}$$

where $m_{3/2}$ is a gravitino mass and a_B is a complex numerical factor whose amplitude is of order of unity. $V_{\not{S}}$ 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 $W_{\rm nr}$. 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,

$$V_{\rm H} = -c_{\rm H}H^2|\Phi|^2, \tag{13}$$

where $c_{\rm H}$ is a positive numerical factor of order of unity. Thus, the AD field evolves with the effective potential,

$$V_{\rm eff} = V + V_{\rm H}. \tag{14}$$

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

$$|\Phi| \simeq (HM_*^n)^{1/(n+1)}, \tag{15}$$

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 $H_{\rm osc}^2 \simeq |V''_{\rm eff}|$, 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,

$$n_B(t_{\rm osc}) \simeq \beta m_{3/2}(H_{\rm osc}M_*^n)^{2/(n+1)}\sin(n\theta_{\rm inf} + \alpha), \tag{16}$$

where $\theta_{\rm inf}$ and α are the phases of Φ during inflation and the constant a_B in the third term of Eq. (12), respectively. Just after the onset of the AD field oscillation, $a^3 n_B$ 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

$$\left(\frac{n_B}{s}\right)_0 \simeq \frac{\beta m_{3/2}T_R}{M_G^2 H_{\rm osc}^2}(H_{\rm osc}M_*^n)^{2/(n+1)}\sin(n\theta_{\rm inf} + \alpha). \tag{17}$$

The Hubble parameter at the onset of the AD field oscillation is

$$H_{\rm osc} \simeq m_0(|\Phi|_{\rm osc}), \tag{18}$$

where $m_0(|\Phi|) \equiv V''_{\not{S}}(|\Phi|)$. Hence, from Eq. (17) we can find that the present value of the baryon-to-entropy ratio is proportional to the reheating temperature.

4 Result

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

$$S_B \equiv \frac{\delta n_B}{n_B} - \frac{\delta s}{s} = \frac{\delta(n_B/s)}{n_B/s} \ . \tag{19}$$

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. From Eq. (**??**), we have

$$S_B = \frac{\delta(n_B/s)}{n_B/s} = \frac{\delta T_R}{T_R} \ . \tag{20}$$

Hence, in the case where the curvature perturbation originates mainly from the modulated reheating mechanism, from Eq. (5) and Eq. (20), we have

$$S_B = -3\zeta \ . \tag{21}$$

The current observational limit for the anti-correlated baryonic-isocurvature fluctuation is roughly given by $|S_B/\zeta| \lesssim O(0.1)$ and hence it means that this model seems to be conflict with current observations. This result is quite similar to that in the case where one consider the gravitino dark matter in the modulated reheating scenario [3]. However, in Ref. [7], we have shown that as for the AD baryogenesis scenario if one consider the thermal effect even before the reheating one can construct the successful AD baryogenesis scenario in modulated reheating scenario.

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