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Affleck-Dine baryogenesis after D-term inflation and solutions to the baryon-dark matter coincidence problem

Masahiro Kawasaki and Masaki Yamada

Institute for Cosmic Ray Research, The University of Tokyo, 5-1-5 Kashiwanoha, Kashiwa, Chiba 277-8582, Japan and Kavli IPMU (WPI), TODIAS, The University of Tokyo, 5-1-5 Kashiwanoha, Kashiwa 277-8583, Japan

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We investigate the Affleck-Dine baryogenesis after D-term inflation with a positive Hubble-induced mass term for a $B-L$ flat direction. It stays at a large field value during D-term inflation, and just after inflation ends it starts to oscillate around the origin of the potential due to the positive Hubble-induced mass term. The phase direction is kicked by higher-dimensional Kähler potentials to generate the $B-L$ asymmetry. The scenario predicts nonzero baryonic isocurvature perturbations, which would be detected by future observations of CMB fluctuations. We also provide a D-term inflation model which naturally explains the coincidence of the energy density of baryon and dark matter.

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I. INTRODUCTION

Inflation is a new paradigm to solve cosmological problems related to the initial conditions of the early Universe. However, any preexisting baryon asymmetry is diluted away due to the inflationary expansion of the Universe, so there should be a mechanism to generate the observed baryon asymmetry after inflation. In addition, the observed abundance of baryon asymmetry is equal to that of dark matter (DM) within order unity ($\Omega_b/\Omega_{\rm DM} \simeq 0.2$) [1], which is a mystery in cosmology referred to as the baryon-DM coincidence problem. This implies that the baryon asymmetry and DM have a common origin. We therefore need to consider inflation, baryogenesis, and DM production simultaneously to construct consistent cosmological models.

In this paper, we focus on D-term inflation models for the following reasons [2,3]: First, the energy scale of D-term inflation is naturally of the order of the GUT scale, which predicts an amplitude of CMB fluctuations consistent with the observations [4,5]. The second reason is related to the so-called η problem. If inflation is driven by a nonzero F-term potential energy, supergravity effects induce masses of the order of the Hubble parameter to all scalar fields, including the inflaton. However, the Hubble-induced mass for the inflaton spoils the flatness of its potential and results in an $\mathcal{O}(1)$ slow roll parameter $\eta \sim 1$. Since Hubble-induced masses come only from nonzero F-term potential energy, the η problem is absent in the case of D-term inflation, in which inflation is driven by nonzero D-term potential energy.

In supersymmetric (SUSY) theories, baryon asymmetry can be generated by Affleck-Dine baryogenesis using a $B-L$ charged flat direction called an AD field [6,7]. Since the soft mass of the AD field is much smaller than the energy scale of inflation, it can stay at a large vacuum expectation value (VEV) during D-term inflation. After D-term inflation ends, the AD field obtains a Hubble-induced mass from the oscillating inflaton field through the supergravity effects. In the literature, the Hubble-induced mass term was assumed to be negative (i.e., tachyonic) or absent to investigate the dynamics of the AD field and calculate the produced baryon asymmetry [8–11]. In this paper, we investigate the Affleck-Dine baryogenesis in the case that the AD field obtains a positive Hubble-induced mass term after D-term inflation. In this case, the AD field starts to oscillate around the origin of the potential due to the positive Hubble-induced mass term just after inflation. At the same time, the phase direction of the AD field feels nonrenormalizable $B-L$ breaking operators and starts to rotate in the phase space, which results in a generation of $B-L$ asymmetry. Then, the coherent oscillation of the AD field decays and dissipates into the thermal plasma, and the $B-L$ asymmetry is converted to the desired baryon asymmetry through the sphaleron effects [12,13]. We calculate the baryon asymmetry and show that the result can be consistent with the observed amount of baryon asymmetry. Since the AD field fluctuates due to the absence of Hubble-induced mass during D-term inflation, the Affleck-Dine baryogenesis predicts some amount of baryonic isocurvature perturbations [9–11,14,15]. Especially, when the AD field starts to oscillate due to the positive Hubble-induced mass term, its radial direction also has quantum fluctuations and contributes to the baryonic isocurvature perurbations. If one were to consider a high-scale D-term inflation model, the resulting isocurvature perturbations would be detected by future observations of CMB fluctuations.

We also build a D-term inflation model which naturally predicts the baryon-to-DM ratio of order unity. We introduce a shift symmetry for the inflaton superfield to ensure a large initial VEV of the inflaton [16]. We also introduce its linear term in the Kähler potential so that the branching of inflaton decay into gravitinos can be of order unity [17–23]. When the mass of a gravitino is larger than $O(100)$ TeV, it decays into the minimal SUSY standard model (MSSM) particles before the big bang nucleosynthesis (BBN) epoch. The DM, which is the lightest SUSY particle (LSP), is therefore produced nonthermally from the gravitino decay. This scenario, together with the above Affleck-Dine baryogenesis scenario, predicts an $\mathcal{O}(1)$ ratio of the energy density of baryon and DM. This means that the scenario naturally explains the baryon-DM coincidence problem. This arises from the fact that both of them are related to the energy scale of inflation. The amount of baryon asymmetry is proportional to the reheating temperature of the Universe and inversely proportional to the Hubble parameter during inflation. The DM abundance is proportional to the reheating temperature and inversely proportional to the mass of the inflaton. Since the Hubble parameter and the mass are related to each other, the resulting baryon and DM density is naturally of order unity. We predict that the LSP mass is 2 orders of magnitude larger than the proton mass, which comes from the fact that the GUT scale is 2 orders of magnitude less than the Planck scale. When the LSP is mostly wino or Higgsino, it would be detected by future indirect detection experiments of DM.

This paper is organized as follows: In the next section, we briefly review the simplest D-term inflation model as an illustration. In Sec. III, we consider the Affleck-Dine baryogenesis and calculate the baryon asymmetry and baryonic isocurvature fluctuations in the case that the AD field obtains a positive Hubble-induced mass term after inflation. In Sec. IV, we provide a D-term inflation model which naturally predicts the baryon-to-DM ratio of order unity. Section V is devoted to the summary.

II. D-TERM INFLATION

We focus on D-term inflation [2,3], in which inflation is driven by a finite energy density of the D-term potential. Although the following simple model of D-term inflation predicts the spectral index relatively blue tilted compared with the observation of CMB fluctuations, we review it as an illustration. Note that there are variants of D-term inflation models which predict a spectral index consistent with the observed value [24,25], and the results in the next section can be applied to those models, too.[1]

We introduce a $U(1)$ gauge symmetry with a Fayet-Ilipoulos (FI) term ξ and consider superfields S, ψ_-, and ψ_+ with $U(1)$ gauge charges as 0, -1, and 1, respectively.[2] The D-term potential is written as

$$V_D = \frac{g^2}{2}(|\psi_+|^2 - |\psi_-|^2 - \xi)^2, \tag{1}$$

where g is the $U(1)$ gauge coupling constant. We introduce a superpotential given as

$$W^{(\mathrm{inf})} = \lambda S \psi_+ \psi_-, \tag{2}$$

where λ is a coupling constant. Hereafter, we denote their scalar components by the same symbols as the superfields.

The scalar component of the field S plays the role of the inflaton. Suppose that the inflaton S has a VEV larger than the critical value of $S_c \equiv g\sqrt{\xi}/\lambda$. The fields ψ_- and ψ_+ obtain large effective masses from the VEV of the inflaton and stay at the origin of the potential. In this regime, the nonzero D-term potential of $V_0 = g^2\xi^2/2$ drives inflation. The Coleman-Weinberg potential for the inflaton lifts its potential above the critical point as

$$V_{\text{1-loop}} \simeq \frac{1}{2} g^2 \xi^2 \left(1 + \frac{g^2}{16\pi^2} \log \frac{\lambda^2 |S|^2}{Q^2}\right), \tag{3}$$

where Q is a renormalization scale. Thus, the inflaton slowly rolls down to the origin of the potential. The COBE normalization requires [4,5]

$$\sqrt{\xi} \simeq 6.6 \times 10^{15}\ \mathrm{GeV}. \tag{4}$$

This leads to a Hubble parameter during inflation according to

$$H_I \simeq \frac{g\xi/\sqrt{2}}{\sqrt{3} M_{\mathrm{Pl}}} \simeq 3.7 \times 10^{12}\ \mathrm{GeV} \left(\frac{\sqrt{\xi}}{6.6 \times 10^{15}\ \mathrm{GeV}}\right)^2. \tag{5}$$

The e-folding number and a slow roll parameter are calculated as

$$N_* \simeq \frac{4\pi^2}{g^2} \frac{S_*^2}{M_{\mathrm{Pl}}^2}, \tag{6}$$

$$\eta \equiv \frac{V''}{V} M_{\mathrm{Pl}}^2 \simeq -\frac{g^2}{8\pi^2} \frac{M_{\mathrm{Pl}}^2}{S_*^2} \simeq -\frac{1}{2N_*}, \tag{7}$$

where the subscript $_*$ denotes values corresponding to the pivot scale $k_* = 0.05\ \mathrm{Mpc}^{-1}$. The slow roll condition fails ($\eta \sim 1$) at the VEV around $S \simeq g/(2\sqrt{2}\pi)$, which is larger than the critical value S_c for the case of $\lambda = \mathcal{O}(1)$. Thus, slow roll inflation ends at the VEV around $S \simeq g/(2\sqrt{2}\pi)$, and soon after that, the waterfall field ψ_+ starts to oscillate around the low-energy minimum of $\sqrt{\xi}$.

The scalar spectral index is calculated as

$$n_s \simeq \frac{1}{N_*} \simeq 0.98. \tag{8}$$

[1]The results are not applicable to the inflation model considered in Ref. [26], because an F-term potential drives inflation with a sizable e-folding number in that scenario.

[2]The FI term can be generated dynamically as discussed in Ref. [27]. They also argue that the production of cosmic strings at the end of inflation is generally avoided when the FI term is generated dynamically. Otherwise, the CMB data puts an upper bound on a cosmic string contribution to the CMB fluctuations, which leads to the upper bound on the FI term [24,28].

It is measured by Planck data alone, such as [5]

$$n_s^{(\mathrm{obs})} = 0.9616 \pm 0.0094 \quad (68\%). \tag{9}$$

So, the above prediction deviates by about 2σ from the observation. Let us emphasize that the results in the next section can be applied to other variants of D-term inflation models, including the ones which predict a spectral index consistent with the observed value within the 1σ level [24,25].

After inflation ends, the energy density of the Universe is dominated by that of the oscillation of S and ψ_+. When some MSSM fields carry nonzero $U(1)$ charge, the field ψ_+ immediately decays into the MSSM fields through the interaction in the D-term potential. Even if the MSSM fields have no $U(1)$ charge, the kinetic mixings between the $U(1)$ and $U(1)_Y$ make the field ψ_+ decay into the MSSM fields relatively fast [8]. Thus, the reheating temperature of the Universe is determined by the relatively late-time decay of the inflaton S, which dilutes the relics produced from the decay of ψ_+. We define the reheating temperature as

$$T_{\mathrm{RH}} \simeq \left(\frac{90}{g_*(T_{\mathrm{RH}})\pi^2}\right)^{1/4} \sqrt{\Gamma_S M_{\mathrm{Pl}}}, \tag{10}$$

where Γ_S is the decay rate of the inflaton S. The reheating temperature of the Universe depends on the mass of the inflaton ($m_S \equiv \lambda\sqrt{\xi}$) and the assumption of interactions between S and the MSSM fields, which is determined by their underlying symmetry. We explicitly calculate the reheating temperature in Sec. IV for a specific model with a shift symmetry and an approximate Z_2 symmetry, while we regard it as a free parameter in the next section to calculate the amount of baryon asymmetry generated by Affleck-Dine baryogenesis.

III. AFFLECK-DINE BARYOGENESIS

In this section, we consider the Affleck-Dine baryogenesis [6,7] after D-term inflation and calculate the resulting baryon asymmetry and baryonic isocurvature perturbations. We consider the case in whih the AD field obtains a positive Hubble-induced mass term after the end of inflation, which has been overlooked in the literature. We investigate the potential of the AD field after D-term inflation in the next subsection, and then we calculate the baryon asymmetry. In Sec. III C, we calculate baryonic isocurvature perturbations predicted by the Affleck-Dine baryogenesis. Finally, we comment on Q-ball formation in Sec. III D.

A. Potential of the AD field

We consider the Affleck-Dine baryogenesis [6,7] using a flat direction (= AD field) ϕ with nonzero $B-L$ charge. The AD field has soft SUSY-breaking terms through the low-energy SUSY-breaking effect. Since the soft mass of the flat direction is much smaller than the Hubble parameter during inflation, it has a large VEV during inflation. In this paper, we assume that the superpotential of the AD field is absent or sufficiently small so that the initial VEV of the AD field ϕ_i can be as large as the Planck scale ($\phi_i \simeq M_{\mathrm{Pl}}$).[3] Such a large VEV is favored to avoid the baryonic isocurvature constraint as shown in Sec. III C (see also Ref. [15]). Note that owing to the exponential term in the supergravity potential, the VEV of the AD field is restricted below the Planck scale. Since the curvature of the phase direction is absent (or at least much less than the Hubble parameter), the phase of the flat direction also stays at a certain phase during inflation. We denote the initial phase of the AD field as θ_i.

After inflation, the energy density of the Universe is dominated by that of oscillating inflatons. Since the inflaton oscillation induces F-term potential, the flat direction obtains Hubble-induced terms through supergravity effects [7]. The scalar potential in supergravity is given as

$$V = e^{K/M_{\mathrm{Pl}}^2}\left[(D_i W)K^{i\bar{j}}(D_j W)^* - \frac{3}{M_{\mathrm{Pl}}^2}|W|^2\right], \tag{11}$$

where K is a Kähler potential and $D_i W \equiv W_i + K_i W/M_{\mathrm{Pl}}^2$. The subscripts represent the derivative with respect to field i and $K^{i\bar{j}} \equiv (K_{i\bar{j}})^{-1}$. Since the F-term of ψ_- is given by $\lambda S\psi_+$, the scalar potential includes the following term:

$$V \supset \frac{|\phi|^2}{M_{\mathrm{Pl}}^2}|\lambda S\psi_+|^2. \tag{12}$$

Using $\psi_+ = \sqrt{\xi}$ and taking the average with respect to time, we obtain the Hubble-induced mass term of

$$V \supset \frac{3}{2}H^2(t)|\phi|^2. \tag{13}$$

Here we have used the virial theorem:

$$m_S^2\langle|S|^2\rangle \simeq \frac{3}{2}H^2(t)M_{\mathrm{Pl}}^2, \tag{14}$$

where $\langle\rangle$ represents the time average. In the literature, they assume a negative Hubble-induced mass term which comes from higher-dimensional Kähler potentials, say,

$$K^{(H)} = c_S|S|^2\frac{|\phi|^2}{M_{\mathrm{Pl}}^2}, \tag{15}$$

[3] For example, we can introduce R symmetry, in which the charge of the AD field is zero to forbid the superpotential for the AD field. This symmetry is not exact, because it is inconsistent with the constant term in the superpotential, which is needed to ensure the (almost) vanishing cosmological constant as well as the gaugino mass terms. Since the R-symmetry-breaking order parameter is of the order of the gravitino mass, which is much smaller than the energy scale of inflation, such breaking terms can be neglected in the following discussion.

where c_S is an $\mathcal{O}(1)$ constant. As we stated in the Introduction, we consider the positive Hubble-induced mass term and do not need to introduce such higher-dimensional terms. Hereafter, we consider a general case and denote the coefficient of the Hubble-induced mass term as c_H [$= \mathcal{O}(1)$].

In addition to the Hubble-induced mass term, the flat direction obtains higher-dimensional terms from nonrenormalizable Kähler potentials.[4] The following Kähler potential may exist and induce $U(1)$-breaking higher-dimensional potential for the flat direction:

$$\begin{aligned} &\frac{2}{3}a_H \int \mathrm{d}^2\theta \mathrm{d}^2\bar{\theta}|S|^2 \frac{\phi^n}{nM_{\rm Pl}^n} + \text{c.c.} \\ &\simeq -\frac{2}{3}a_H|\partial_\mu S|^2 \frac{\phi^n}{nM_{\rm Pl}^n} + \text{c.c.} \\ &\simeq -a_H H^2(t)\left(\frac{\phi^n}{nM_{\rm Pl}^{n-2}} + \text{c.c.}\right), \end{aligned} \tag{16}$$

where n is an integer depending on the flat directions. For example, $n = 3, 6, 9, \ldots$ for the $u^c d^c d^c$ flat direction. In the last line, we take the average with respect to time and use the relation of $\langle(\partial_0 S)^2\rangle \simeq 3H^2(t)M_{\rm Pl}/2$, which comes from the virial theorem. Note that this term has a nonzero phase which is different from the phase of the flat direction θ_i during inflation. We can redefine the phase of the flat direction to eliminate the phase of a_H. After the elimination, we redefine the initial phase of the flat direction as θ_i without loss of generality. The discrepancy between the initial phase of the flat direction and the phase of the above $U(1)$ breaking term is essential to generate the baryon asymmetry.

In summary, the AD field obtains the following potential after inflation:

$$V(\phi) = c_H H^2(t)|\phi|^2 - a_H H^2(t)\left(\frac{\phi^n}{nM_{\rm Pl}^{n-2}} + \text{c.c.}\right) + \cdots, \tag{17}$$

where c_H and a_H are positive $\mathcal{O}(1)$ parameters. The dots represents higher-dimensional terms which restrict the AD field below the Planck scale. Since the flat direction starts to oscillate due to the Hubble-induced mass, we can neglect usual soft mass and A-terms for the AD field.

B. Calculation of baryon asymmetry

In this subsection, we calculate the baryon asymmetry generated from the AD field with the potential (17). The initial VEV and phase are ϕ_i and θ_i, respectively. For $c_H > 0$, the flat direction starts to oscillate around the origin of the potential just after the end of inflation. At the same time, the flat direction is kicked in the phase direction due to the second term in Eq. (17). The $B - L$ asymmetry is generated through this dynamics. The evolution of the equation for the $B - L$ charge density is written as

$$\dot{n}_{B-L} + 3Hn_{B-L} = 2q\,\mathrm{Im}\left[\phi^* \frac{\partial V}{\partial \phi^*}\right], \tag{18}$$

where q denotes the $B - L$ charge of the AD field. From this equation we obtain

$$\begin{aligned} a^3 n_{B-L}(t_{\rm osc}) &\simeq \int \mathrm{d}t 2qa^3(t)|\phi V_A'| \sin(n\theta) \\ &\equiv \epsilon q H_I \phi_i^2, \end{aligned} \tag{19}$$

$$\epsilon \simeq (3-4) \times \frac{8}{3n-6} a_H \sin(-n\theta_i)\left(\frac{\phi_i}{M_{\rm Pl}}\right)^{n-2}, \tag{20}$$

where we have used $\phi \propto a^{-3/4}$. We define ϵ (≤ 1), which represents the efficiency of baryogenesis. We have numerically solved the equations of motion for ϕ and S with the Friedmann equation and have obtained the numerical factor of (3–4) for $\epsilon \lesssim 1$. Since the baryon density has to be smaller than the number density of the AD field, ϵ is at most unity even for large a_H and ϕ_i. The amplitude of the flat direction decreases as time evolves due to the Hubble expansion, and the $B - L$ breaking effect is absent soon after the oscillation. Thus, the generated $B - L$ asymmetry is conserved soon after the AD field starts to oscillate.

Then, the oscillating AD field decays and dissipates into radiation [29], and the sphaleron effect converts the $B - L$ asymmetry to the baryon asymmetry [12,13]. Since the sphaleron process is in thermal equilibrium, the resultant baryon asymmetry is related to the $B - L$ asymmetry as [30]

$$n_b = \frac{8}{23} n_{B-L}. \tag{21}$$

Assuming the absence of entropy production other than the reheating by inflaton decay, we can calculate the resulting baryon-to-entropy ratio Y_b as

$$\begin{aligned} Y_b &\equiv \frac{n_b}{s} = \frac{8}{23}\frac{n_{B-L}}{s}\bigg|_{\rm RH} \simeq \frac{8}{23}\frac{3T_{\rm RH}n_{B-L}}{4\rho_S}\bigg|_{\rm osc} \simeq \frac{8}{23}\frac{\epsilon q T_{\rm RH}}{4H_I}\left(\frac{\phi_i}{M_{\rm Pl}}\right)^2 \\ &\simeq 8.7 \times 10^{-11} \epsilon q \left(\frac{T_{\rm RH}}{4 \times 10^3\ \text{GeV}}\right)\left(\frac{H_I}{4 \times 10^{12}\ \text{GeV}}\right)^{-1}\left(\frac{\phi_i}{M_{\rm Pl}}\right)^2, \end{aligned} \tag{22}$$

[4] The usual Hubble-induced A-terms are absent during the inflaton oscillation era.

where ρ_S ($\simeq 3H_I^2 M_{\rm Pl}^2$) is the energy density of the inflaton S. This can be consistent with the observed baryon asymmetry of $Y_b^{\rm obs} \simeq 8.7 \times 10^{-11}$ [1].

One can find differences from the conventional scenario of the Affleck-Dine baryogenesis. The Hubble parameter at the beginning of oscillation H_I is determined by the energy scale of inflation, not by the curvature of the potential for the flat direction (see Ref. [15], for example). This is because the flat direction starts to oscillate just after the end of inflation due to the positive Hubble-induced mass term, while in the conventional scenario it starts to oscillate at $H(t) \simeq m_\phi$, where m_ϕ is the soft mass of the AD field. This allows us to consider a relatively large reheating temperature even if the initial VEV ϕ_i is as large as the Planck scale. In addition, the ellipticity parameter ϵ, which describes the efficiency of baryogenesis, can be much smaller than unity when ϕ_i is smaller than the Planck scale. This is because the phase direction of the AD field is kicked by a higher-dimensional Kähler potential, which is highly suppressed for a small VEV of the AD field. However, as shown in the next subsection, the baryonic isocurvature constraint requires the initial VEV to be as large as the Planck scale. In that case, ϵ is of order unity for $a_H = \mathcal{O}(1)$.

One might wonder if the energy density of the AD field dominates that of the Universe in the case in which its initial VEV is as large as the Planck scale. This may be true in the case of conventional Affleck-Dine baryogenesis, in which the AD field starts to oscillate when the Hubble parameter decreases down to the soft mass of the flat direction. However, the energy density of the AD field never dominates that of the Universe in the above scenario, because it decreases faster than that of radiation. Just after inflation, the AD field starts to oscillate around the origin due to the positive Hubble-induced mass term. Then, its number density decreases with time as a^{-3} due to the expansion of the Universe. This means that its energy density decreases as $a^{-9/2}$ because its effective mass is of the order of the Hubble parameter, which decreases as $a^{-3/2}$. When the Hubble parameter decreases down to the mass of the AD field [that is, when $H(t) \simeq m_\phi$], its energy fraction compared to the total energy density is given as

$$\left.\frac{\rho_{\rm AD}}{\rho_{\rm tot}}\right|_{H\simeq m_\phi} \simeq \left(\frac{m_\phi}{H_I}\right)\left.\frac{\rho_{\rm AD}}{\rho_{\rm tot}}\right|_{H\simeq H_I} \simeq 10^{-11}\left(\frac{\phi_i}{M_{\rm Pl}}\right)^2\left(\frac{m_\phi}{\rm TeV}\right)\left(\frac{H_I}{4\times 10^{12}\ {\rm GeV}}\right)^{-1}. \tag{23}$$

Thus, the energy density of AD field becomes negligible soon after inflation, and the result of Eq. (22) is applicable to the case of $\phi_i \simeq M_{\rm Pl}$.

C. Baryonic isocurvature perturbations

Although the initial phase and radial values of the flat direction are almost constant over the whole range of the observable Universe, they acquire quantum fluctuations like [9–11,14,15]

$$|\delta\theta_i| \simeq \frac{H_I}{2\pi|\phi_i|}, \tag{24}$$

$$|\delta\phi_i| \simeq \frac{H_I}{2\pi}. \tag{25}$$

These fluctuations result in baryonic isocurvature perturbations because the produced baryon density is related to the initial phase θ_i and VEV ϕ_i [see Eqs. (19) and (20)]. The baryonic isocurvature perturbation $S_{b\gamma}$ is given by

$$\mathcal{S}_{b\gamma} \equiv \frac{\delta Y_B}{Y_B} \simeq n\left(\cot(n\theta)\delta\theta + \frac{\delta\phi_i}{\phi_i}\right). \tag{26}$$

Since we consider the case in which the AD field starts to oscillate due to the positive Hubble-induced mass term, its radial direction also has quantum fluctuations and contributes to the baryonic isocurvature perturbations. This leads to an additional factor in Eq. (26), which cannot be suppressed by the tuning of the initial phase θ_i. Note that the VEV of the AD field is smaller than the Planck scale due to the exponential factor in the supergravity potential. This leads to a lower bound on baryonic isocurvature perturbations like

$$|S_{{\rm b}\gamma}| \simeq 2.7\times 10^{-7}\times n\left(\frac{H_I}{4\times 10^{12}\ {\rm GeV}}\right)\left(\frac{M_{\rm Pl}}{\phi_i}\right), \gtrsim 2.7\times 10^{-7}\times n\left(\frac{H_I}{4\times 10^{12}\ {\rm GeV}}\right), \tag{27}$$

where we assume $1+\cot(n\theta) \approx 1$.

Since the density perturbations of the cosmic microwave background are predominantly adiabatic [4,5], the baryonic isocurvature perturbation is tightly constrained as [15,31]

$$|S_{{\rm b}\gamma}| \lesssim 5.0\times 10^{-5}. \tag{28}$$

Our scenario predicts a value below this constraint, though it depends on the value of n and H_I (i.e., ξ). If one might consider a high-scale D-term inflation model, the resulting isocurvature perturbations can be as large as this lower bound and would be detected by future observations of CMB fluctuations.

D. Comments on Q-ball formation

In this subsection, we comment on Q-ball formation. If the potential of the AD field is shallower than the quadratic potential, its coherent oscillation is unstable and fragments into nontopological solitons, called Q-balls [32]. The formation of Q-balls may change the scenario of Affleck-Dine baryogenesis [33–39]. For example, their decay can be another source of nonthermal production of DM [15,35,40–51], or Q-balls can be a candidate for DM if they are stable [34,52–54]. In the case considered in this paper, the AD field starts to oscillate by the positive Hubble-induced mass term. When the beta function for the Hubble-induced mass of the AD field is positive, Q-balls do not form. The beta function has positive contributions from Yukawa interactions, while it has negative ones from gauge interactions. The former positive contributions are roughly proportional to the squared masses of squarks and sleptons, and the latter negative ones are roughly proportional to the squared masses of gauginos. Here, since the Hubble-induced mass for gauginos is absent or one-loop suppressed, the positive contributions from Yukawa interactions are usually dominant. Therefore, the beta function for the Hubble-induced mass of the AD field is usually positive, and Q-balls may not form in our scenario. However, if the AD field consists only of the first- and second-family squarks and/or sleptons, the positive contributions from Yukawa interactions are suppressed by small Yukawa couplings. In this case, Q-balls might form. We estimate the typical charge Q of Q-balls as

$$Q \sim \beta \left(\frac{\phi_i}{m_{\phi,\mathrm{eff}}} \right)^2, \tag{29}$$

where $m_{\phi,\mathrm{eff}}$ is the effective mass and β ($\sim 10^{-2}$) is a numerical factor obtained from simulations of Q-ball formation [36,37,55]. Here, we should substitute the Hubble-induced mass into the effective mass, and so the typical charge of Q-balls is at most 10^8. Such small Q-balls soon evaporate into thermal plasma via interactions with the thermal plasma [56,57] (see also Ref. [38]).[5] Therefore, the subsequent cosmological scenario and the calculation of the baryon asymmetry does not change.

Even if Q-balls do not form just after the end of inflation, they may form at the time of $H(t) \simeq m_\phi$. After that time, the potential of the AD field is dominated by its soft mass term. If the beta function of the soft mass is negative, the AD field begins to fragment into Q-balls at that time. Since $n_b \propto H(t)\phi^2(t) \propto a^{-3} \propto H(t)^2$ until the Hubble parameter decreases down to the soft mass, the amplitude of the AD field at $H(t) \simeq m_\phi$ is given as

$$\phi|_{H(t)\simeq m_\phi} \simeq \left(\frac{m_\phi}{H_I} \right)^{1/2} \phi_i. \tag{30}$$

[5] The evaporation is efficient during the inflaton oscillation era. In addition, since the energy per unit charge for these Q-balls is given by the Hubble parameter, their energy density decreases with time as $a^{-9/2}$. Thus, the energy density of the Q-balls never dominates that of the Universe.

This implies that a typical charge of Q-balls is given as

$$Q \simeq \beta \left(\frac{\phi}{m_\phi} \right)^2 \simeq \beta \left(\frac{\phi_i^2}{m_\phi H_I} \right). \tag{31}$$

This is at most 10^{18} for typical parameters. Such small Q-balls evaporate into thermal plasma soon after they form. Even if Q-balls survive, they are so small as to decay into quarks before the BBN epoch. However, they usually decay after the electroweak phase transition [58,59]. Since the sphaleron process is decoupled at that time, the AD field has to carry a nonzero baryon charge (not $B-L$) to generate the baryon asymmetry. In that case, the resulting baryon asymmetry is given by Eq. (22) without the factor of $8/23$.

IV. MODEL FOR SOLUTION TO THE BARYON-DM COINCIDENCE PROBLEM

In this section, we propose a D-term inflation model which predicts an $\mathcal{O}(1)$ ratio of baryon to DM density. We introduce a shift symmetry to ensure the flatness of the inflaton potential above the Planck scale. When there exists a small linear term in the Kähler potential, the inflaton decays mainly into gravitinos [17–23]. The subsequent decay of those gravitinos is a source of nonthermal production of LSP DM, and the resulting DM abundance is proportional to the reheating temperature and inversely proportional to the inflaton mass. Since the amount of the baryon asymmetry in Eq. (22) has similar parameter dependences, the baryon and DM densities are related with each other through the energy scale of inflation.

In the next subsection, we propose the D-term inflation model. In Sec. IV B, we investigate reheating processes of the D-term inflation model and then calculate the abundance of DM.

A. Model

Let us introduce a shift symmetry and an approximate Z_2 symmetry for the inflaton field S [16]. Under these symmetries, S transforms as $S \to S + i\alpha$ (α:real) and $S \to -S$, respectively. Then, the Kähler potential is written as

$$K = c_S(S + S^*) + \frac{1}{2}(S + S^*)^2 + |\psi_-|^2 + |\psi_+|^2, \tag{32}$$

where c_S ($\ll 1$) is an order parameter for the Z_2 symmetry breaking effect. The superpotential of Eq. (2) explicitly breaks the shift symmetry, which is required to ensure a graceful exit. Otherwise, the inflaton stays at a certain VEV because it has an exactly flat potential. In this model, we should replace $|S|^2$ with $(S + S^*)^2/2$ for the calculations in Sec. IV A, though the results are unchanged.

There is an advantage to imposing the shift symmetry on the inflaton. In order to obtain a sufficiently large e-folding number, say, $N_* \gtrsim 60$, the initial VEV of the inflaton S has to be as large as $N_* \frac{\sqrt{2}g^2}{4\pi^2} M_{\rm Pl} \simeq 0.5 M_{\rm Pl}$. which is of the order of the Planck scale. This implies that the Planck-scale physics may affect the potential of the inflaton and spoil its flatness. However, the shift symmetry ensures the flatness of the inflaton potential above the Planck scale.

If Z_2 symmetry is exact, some MSSM particles have to carry odd Z_2 charge and interact with the field S [16] for the inflaton to decay. In this case the LSP DM abundance is given by the usual thermal relic density. Here, we introduce Z_2 breaking terms in the Kähler potential so that the field S efficiently decays into gravitinos [22,23], whose decay is a source of nonthermal production of LSP DM.

We assume that the mass of gravitinos is of order $10^{2\text{–}3}$ TeV so that gravitinos decay into radiation before the BBN epoch. Otherwise, the decay of gravitinos spoils the success of the BBN, or their energy density overcloses the Universe if they are stable. Such a heavy gravitino is well motivated in a class of SUSY models with a split spectrum [60–66]. In these models, the masses of gravitinos as well as squarks and sleptons are of order (or larger than) $10^{2\text{–}3}$ TeV, while those of gauginos are of order 1 TeV. This hierarchy can be realized when gauginos acquire one-loop suppressed soft masses through the anomaly-mediated SUSY-breaking effect [67,68]. In that case, the masses of the wino and gravitino are related to each other such as

$$m_{\tilde{w}} = \frac{g_2^2}{16\pi^2}(m_{3/2} + L) \simeq 3 \times 10^{-3}(m_{3/2} + L). \tag{33}$$

The factor L is the Higgsino threshold correction and is calculated as [68,69]

$$L \equiv \mu_H \sin 2\beta \frac{m_A^2}{|\mu_H|^2 - m_A^2} \log \frac{|\mu_H|^2}{m_A^2}, \tag{34}$$

where m_A is the mass of the heavy Higgs bosons, μ_H is the SUSY mass of the Higgsinos, and $\tan\beta$ is the ratio of the VEVs of H_u and H_d. When μ_H is of the order of the gravitino mass, the Higgsino threshold correction is important, and the wino mass is $\sim 10^{-3} m_{3/2}$. Note that a neutral Higgsino can also be the LSP when μ_H is sufficiently small. The following discussion does not rely on the detailed properties of the LSP except for its mass. Hereafter, we assume that the mass of the gravitino is $O(10^{2\text{–}3})$ TeV and that of the LSP is $O(10^{2\text{–}3})$ GeV.

In order to calculate the gravitino production rate from the inflaton decay, we need to specify the SUSY-breaking sector. We introduce a Polonyi field z, which breaks SUSY in the low-energy scale, and consider a simple extension of the Polonyi model given as

$$K = |z|^2 - \frac{|z|^4}{\Lambda^2}, \tag{35}$$

$$W = \mu^2 z + W_0, \tag{36}$$

where Λ is a cutoff scale, μ is the SUSY-breaking scale, and W_0 is a constant term which makes the cosmological constant (almost) zero in the present Universe. This can be achieved by the O'Raifeartaigh model after integrating out relatively heavy particles [70] or by dynamical SUSY breaking models, including the IYIT model [71,72]. The important parameters are calculated as

$$\mu^2 \simeq \sqrt{3} m_{3/2} M_{\rm Pl}, \tag{37}$$

$$m_z^2 \simeq \frac{12 m_{3/2}^2}{\Lambda^2} M_{\rm Pl}^2, \tag{38}$$

$$\langle z \rangle_0 \simeq 2\sqrt{3}\left(\frac{m_{3/2}}{m_z}\right)^2 M_{\rm Pl}, \tag{39}$$

where m_z is the mass of z and $\langle z \rangle$ is its VEV at the low-energy vacuum. Since the Hubble-induced mass is absent during D-term inflation, massless scalar fields cannot be stabilized at the origin. This implies that if the mass of the Polonyi field is much smaller than the Hubble parameter, it obtains a VEV as large as the Planck scale during inflation. In this section, we consider the case that the mass of the Polonyi field is larger than H_I and it stays at the origin of the potential during inflation, while we consider the case of relatively light Polonyi in the Appendix. The conditions of $m_z \gtrsim H_I$ and $\Lambda \gtrsim \mu$ in the effective theory lead to the lower bound on the gravitino mass [73]:

$$\begin{aligned} m_{3/2} &\gtrsim \frac{\sqrt{3} H_I^2}{12 M_{\rm Pl}}, \\ &\simeq 10^3 \text{ TeV}\left(\frac{H_I}{4 \times 10^{12} \text{ GeV}}\right)^2. \end{aligned} \tag{40}$$

Thus, the heavy gravitino is favored in D-term inflation to stabilize the VEV of the Polonyi field.

B. Reheating process

In this subsection, we investigate the reheating process and calculate the DM abundance in the model introduced in the previous subsection.

As explained in Sec. II, the field ψ_+ decays into the MSSM fields much faster than the inflaton S, so that the reheating temperature of the Universe is determined by the relatively late-time decay of the inflaton S [8]. After the field ψ_+ decays completely, the effective superpotential can be rewritten as

$$W^{(\rm inf)} = m_S S \psi_-, \tag{41}$$

$$m_S \equiv \lambda\sqrt{\xi}, \tag{42}$$

where m_S is the effective mass of the fields S and ψ_-. This superpotential is equivalent to that in the model of chaotic inflation proposed in Ref. [16], except for the value of m_S.

The supergravity effects induce a soft SUSY-breaking B-term of $bm_{3/2}m_S S\psi_-$, where b is an $\mathcal{O}(1)$ constant. This implies that they maximally mix with each other and form mass eigenstates

$$\Phi_\pm \equiv \frac{1}{\sqrt{2}}(S \pm \psi_-^\dagger) \tag{43}$$

around the potential minimum [18,22]. Therefore, when the time scale of inflaton decay Γ_S^{-1} is longer than that of the mixing effect $m_{3/2}^{-1}$, we have to consider the decay of $\Phi_\pm$ to investigate the reheating process. Since we consider a heavy gravitino and a reheating temperature of order 10^3 GeV [see Eq. (22)], the mixing effect is indeed relevant.

The Z_2 breaking term in the Kähler potential results in the decay of the field $\Phi_\pm$ through supergravity effects [23].[6] First, let us focus on the top Yukawa interaction in the MSSM sector:

$$W^{(\mathrm{top})} = y_t Q_3 H_u u_3^c, \tag{44}$$

where y_t is the top Yukawa coupling constant, and Q_3, H_u, and u_3^c are the chiral supermultiplets of the MSSM sector. The relevant interaction terms between ψ_- and the MSSM fields are given by

$$\begin{aligned} V &= \frac{1}{M_{\mathrm{Pl}}} K_S W^{(\mathrm{top})} W_S^* + \mathrm{c.c.} + \cdots, \\ &= \frac{y_t m_S K_S}{M_{\mathrm{Pl}}^2} \psi_-^* (\tilde{Q}_3 H_u \tilde{u}_3^c) + \mathrm{c.c.} + \cdots, \end{aligned} \tag{45}$$

where the dots "$\cdots$" represent the other irrelevant terms. Since the fields $\Phi_\pm$ consist of ψ_- as Eq. (43), they decay into the MSSM scalar fields through this interaction. They also decay into the MSSM fermion fields, which equally contributes to the $\Phi_\pm$ decay [23]. Thus, the partial decay rate of $\Phi_\pm$ into the MSSM fields is given as

$$\Gamma_{\mathrm{MSSM}}(\Phi_\pm \to \mathrm{MSSM}) = \frac{3c_S^2}{256\pi^3}|y_t|^2 \frac{m_S^3}{M_{\mathrm{Pl}}^2}, \tag{46}$$

where we use $K_S = c_S$. Since we consider the gaugino mass ($m_{\tilde{g}}$) much smaller than the gravitino mass, the decay rates of S into gauge fields are suppressed by a factor of $(m_{\tilde{g}}/m_{3/2})^2$ and can be neglected [21].

Next, let us consider the decay of $\Phi_\pm$ into gravitinos [17–23]. We follow the discussion presented in Ref. [74]. When the field S has a nonzero VEV, the field ψ_- mixes with the SUSY-breaking field z and can decay into a Goldstino, (i.e., the longitudinal component of a gravitino). This is because the supergravity effects induce mixing terms such as

$$\begin{aligned} V &= W_S (K_S W)^* + K_{S\bar{z}}^{-1} W_S W_z^* + \mathrm{c.c.} + \cdots \\ &= m_S d F_z \psi_- z^* + \mathrm{c.c.} + \cdots, \end{aligned} \tag{47}$$

$$d \equiv \langle K_S \rangle - \langle K_{Sz\bar{z}} \rangle, \tag{48}$$

where the dots "$\cdots$" represent the other irrelevant terms. The second term in d is relevant when there is a term like $(S+S^*)|z|^2$ in the Kähler potential, whose coefficient is of order c_S. Thus, the fields ψ_- and z mix with each other, and the mixing angle is given by

$$\theta \simeq d\frac{F_z m_S}{m_z^2}, \tag{49}$$

where we use $m_z \gg m_S$. Since the fields $\Phi_\pm$ consist of ψ_-, they mix with z, and the mixing angle is given by $\theta/\sqrt{2}$. Since the SUSY-breaking field z has an operator of

$$\mathcal{L} = -2\frac{F_z}{\Lambda^2} z\tilde{z}^\dagger \tilde{z}^\dagger + \mathrm{H.c.}, \tag{50}$$

it decays into a Goldstino $\tilde{z}$. Together with the mixings between $\Phi_\pm$ and z, the fields $\Phi_\pm$ decays into Goldstinos through this operator. The partial decay rate of the field $\Phi_\pm$ into a Goldstino is therefore calculated as [75]

$$\begin{aligned} \Gamma_{\tilde{z}}(\Phi_\pm \to \tilde{z}\,\tilde{z}) &\simeq \frac{1}{32\pi}\left(\frac{\theta}{\sqrt{2}}\right)^2 \frac{m_z^4}{|F_z|^2} m_S, \\ &\simeq \frac{d^2}{64\pi}\frac{m_S^3}{M_{\mathrm{Pl}}^2}. \end{aligned} \tag{51}$$

From Eqs. (46) and (51), the total decay rate Γ_S and the branching ratio of the decay of $\Phi_\pm$ into gravitinos $B_{3/2}$ are given by

$$\Gamma_S = \Gamma_{\mathrm{MSSM}}(\Phi_\pm \to \mathrm{MSSM}) + \Gamma_{\tilde{z}}(\Phi_\pm \to \tilde{z}\,\tilde{z}), \tag{52}$$

$$\mathrm{Br}_{3/2} = \frac{d^2}{d^2 + 3|y_t|^2 c^2/(4\pi^2)}. \tag{53}$$

Since $d/c = \mathcal{O}(1)$ and $y_t = \mathcal{O}(1)$, the branching ratio is almost unity. This means that the energy density of the

[6] Since the field ψ_- has a small VEV of order $cm_{3/2}/m_S$ due to the supergravity effects, it can decay into the MSSM fields through the D-term potential. However, we confirm that its partial decay rate is irrelevant due to the suppression factor coming from its small VEV and can be neglected.

Universe is dominated by that of the gravitinos after the fields $\Phi_\pm$ (i.e., the inflaton S) decay completely.[7]

Since the fields $\Phi_\pm$ are much heavier than gravitinos, the produced gravitinos are highly relativistic. The Lorentz factor for the gravitinos at a time $H^{-1}(t)$ is given as

$$\gamma(t) = \left[\left(\frac{m_S}{m_{3/2}}\right)^2 \frac{H(t)}{\Gamma_S} + 1\right]^{1/2} \simeq \frac{m_S}{m_{3/2}} \left(\frac{H(t)}{\Gamma_S}\right)^{1/2}. \quad (54)$$

The gravitinos decay into MSSM particles with a rate of

$$\Gamma_{3/2} \simeq \gamma^{-1}(t) \frac{1}{48\pi} \frac{\sum_i m_{\tilde{X}_i}^5}{m_{3/2}^2 M_{\rm Pl}^2}, \quad (55)$$

where the summation is taken for all MSSM particles $\tilde{X}_i$. Since we consider a SUSY model with relatively light gauginos and relatively heavy squarks and sleptons, we can roughly estimate the numerator as $24 m_{3/2}^5$. This implies that the gravitino decays into radiation at the temperature

$$\begin{aligned} T_{3/2} &\simeq \left(\frac{90}{g_* \pi^2}\right)^{1/4} \sqrt{\Gamma_{3/2} M_{\rm Pl}} \\ &\simeq 1.1\ \text{MeV} \left(\frac{T_{\rm RH}}{4\times 10^3\ \text{GeV}}\right)^{1/3} \left(\frac{m_S}{5\times 10^{15}\ \text{GeV}}\right)^{-1/3} \left(\frac{m_{3/2}}{400\ \text{TeV}}\right)^{4/3}, \end{aligned} \quad (56)$$

where g_* ($\simeq 10.75$) is the effective number of degrees of freedom at the decay time. We require that the mass of gravitino be of order $10^{2\text{–}3}$ TeV or larger so that its decay completes before the BBN epoch; that is, $T_{3/2} \gtrsim 1$ MeV. Otherwise, the decay particles interact with the light elements and spoil the success of the BBN [78–82]. The gravitino decay temperature $T_{3/2}$ is much smaller than the mass of the LSP, so that the decay of the gravitino is a source of its nonthermal production. Since the energy density of the Universe is dominated by that of gravitinos before they decay, the thermal relic density of the LSP is diluted by the entropy production from the gravitino decay. Therefore, the LSP abundance is determined by the nonthermal production from the gravitino decay. The produced number density of the LSPs is equal to that of the gravitinos due to the R-parity conservation. Note that the annihilation of the produced LSP is usually inefficient at such a low temperature.

The Lorentz factor of the gravitino is of order 10^3 for the reference parameters shown in Eq. (56). This implies that the scale factor of the Universe continues to decrease as a^{-4} from the time of reheating by the decay of $\Phi_\pm$. Although the LSPs are relativistic at the time they are produced from gravitino decay, they lose their energy through interactions with the thermal plasma and soon become nonrelativistic particles [83–85]. Therefore, the LSP DM is cold, even though it is produced nonthermally in this scenario.

C. DM density and baryon-DM coincidence

Let us summarize the scenario of nonthermal production of DM. First, the inflaton S (or $\Phi_\pm$) decays into gravitinos as well as the MSSM particles at $H(t) \simeq \Gamma_S$. Then the energy density of the Universe is dominated by the relativistic gravitinos and decreases as a^{-4}. The gravitinos decay into the MSSM particles just before the epoch of the BBN, and the LSP DM is produced nonthermally. Since the thermal relic density of the LSP is diluted by the entropy production of gravitino decay, its abundance is determined by the gravitino decay. Thus, we can estimate the resultant DM abundance as

$$\begin{aligned} Y_{\rm DM} &\equiv \frac{n_{\rm LSP}}{s} \\ &\simeq \left.\frac{n_{3/2}}{s}\right|_{H=\Gamma_{3/2}} \\ &\simeq \left.\frac{3T_{3/2}}{4}\frac{n_{3/2}}{\rho_{3/2}}\right|_{H=\Gamma_{3/2}} \\ &\simeq \left.\frac{3T_{3/2}}{4}\left(\frac{\Gamma_S}{\Gamma_{3/2}}\right)^{1/2}\frac{n_{3/2}}{\rho_{3/2}}\right|_{H=\Gamma_S} \\ &\simeq \left.\frac{3T_{\rm RH}^{(\rm eff)}}{4}\frac{2{\rm Br}_{3/2} n_S}{\rho_S}\right|_{H=\Gamma_S} \\ &\simeq \frac{3T_{\rm RH}^{(\rm eff)}}{2m_S}, \end{aligned} \quad (57)$$

where we have used ${\rm Br}_{3/2} \simeq 1$ in the last line. We define the effective reheating temperature $T_{\rm RH}^{(\rm eff)}$ by Eq. (10) with the replacement of $g_*(T_{\rm RH}) \to g_*(T_{3/2})$ as

$$\begin{aligned} T_{\rm RH}^{(\rm eff)} &\simeq \left(\frac{90}{g_*(T_{3/2})\pi^2}\right)^{1/4} \sqrt{\Gamma_S M_{\rm Pl}} \\ &\simeq 1.5\times 10^3\ \text{GeV} \left(\frac{m_S}{5\times 10^{15}\ \text{GeV}}\right)^{3/2} \left(\frac{d}{10^{-10}}\right). \end{aligned} \quad (58)$$

[7]Note that since we consider relatively a low reheating temperature $\sim 10^{3\text{–}4}$ GeV, we can neglect the thermal production of gravitinos [76,77].

The reheating temperature is adjusted by the Z_2 symmetry order parameter d to obtain a desirable amount of baryon asymmetry from Eq. (22) or DM from Eq. (57).

Here we take into account the baryon asymmetry generated by the Affleck-Dine baryogenesis. Once we replace the reheating temperature T_{RH} with the effective one $T_{\mathrm{RH}}^{(\mathrm{eff})}$ defined by Eq. (58), the resulting baryon asymmetry is still given by Eq. (22) even in this scenario.[8]

Combining Eqs. (22) and (57), we obtain the following simple relation for the baryon-to-DM ratio:

$$\frac{\Omega_b}{\Omega_{\mathrm{DM}}} \simeq \frac{4}{69} \epsilon q \frac{m_p}{m_{\mathrm{LSP}}} \frac{m_S}{H_I}, \tag{59}$$

where we assume $\phi_i \simeq M_{\mathrm{Pl}}$. Substituting benchmark parameters and the proton mass $m_p \simeq 0.938$ GeV, we obtain

$$\begin{aligned}\frac{\Omega_b}{\Omega_{\mathrm{DM}}} &\simeq 0.22\epsilon q \left(\frac{m_{\mathrm{LSP}}}{400\ \mathrm{GeV}}\right)^{-1} \left(\frac{m_S}{6.6\times 10^{15}\ \mathrm{GeV}}\right) \left(\frac{H_I}{4\times 10^{12}\ \mathrm{GeV}}\right)^{-1}, \\ &\simeq 0.12\epsilon q \lambda g^{-1} \left(\frac{m_{\mathrm{LSP}}}{400\ \mathrm{GeV}}\right)^{-1} \left(\frac{\sqrt{\xi}}{6.6\times 10^{15}\ \mathrm{GeV}}\right)^{-1},\end{aligned} \tag{60}$$

which is naturally of order unity and is consistent with the observed value of $\Omega_b^{(\mathrm{obs})}/\Omega_{\mathrm{DM}}^{(\mathrm{obs})} \simeq 0.2$ [1]. The scenario naturally explains the coincidence of their energy density, known as the baryon-DM coincidence problem.[9] This is because both of them are related to the energy scale of inflation. The amount of baryon asymmetry is proportional to the reheating temperature of the Universe and inversely proportional to the Hubble parameter during inflation. That of DM is proportional to the reheating temperature and inversely proportional to the mass of the inflaton. Since the Hubble parameter ($H_I \sim g\xi/M_{\mathrm{Pl}}$) and the inflaton mass ($m_S = \lambda\sqrt{\xi}$) are related to each other, the resulting baryon and DM density is naturally of order unity. Interestingly, the electroweak scale DM mass [$O(10^2)$ GeV] comes from the fact that the GUT scale, which $\sqrt{\xi}$ is expected to be, is 2 orders of magnitude less than the Planck scale.

Although the result has an $\mathcal{O}(1)$ uncertainty coming mainly from λ and ξ, the LSP with a mass of $O(10^{2\text{–}3})$ GeV is favored in our scenario. If the LSP DM is mostly winos or Higgsinos, the indirect detection experiments of DM put lower bounds on DM mass. The wino DM with $m_{\tilde{w}} \leq 390$ GeV and $2.14\ \mathrm{TeV} \leq m_{\tilde{w}} \leq 2.53$ TeV is excluded [87], while the Higgsino DM with $m_{\tilde{h}} \leq 160$ GeV is excluded [88]. Future indirect detection experiments can detect the wino DM with $m_{\tilde{w}} \leq 1.0$ TeV and $1.66\ \mathrm{TeV} \leq m_{\tilde{w}} \leq 2.77$ TeV [87].

V. SUMMARY

We have considered the Affleck-Dine baryogenesis assuming a positive Hubble-induced mass term after D-term inflation. We have calculated the baryon asymmetry and found that the result is consistent with the observed abundance of baryon asymmetry for reheating temperature $T_{\mathrm{RH}} \sim 10^3$ GeV. There are some differences from the conventional scenarios of the Affleck-Dine baryogenesis where the sign of the Hubble-induced mass term is negative. First, since the AD field starts to oscillate just after the end of inflation, the resulting baryon asymmetry is inversely proportional to the inflation scale H_I. This allows us to consider a relatively large reheating temperature even if the initial VEV ϕ_i is as large as the Planck scale. In addition, the ellipticity parameter ϵ, which describes the efficiency of baryogenesis, can be much smaller than unity when ϕ_i is smaller than the Planck scale. This is because the phase direction of the AD field is kicked by a higher-dimensional Kähler potential, which effect is highly suppressed for a small VEV of the AD field. Since the radial direction as well as the phase have quantum fluctuations during D-term inflation, Affleck-Dine baryogenesis predicts nonzero baryonic isocurvature perturbations. They would be detected by future CMB observations if one considers a high-scale D-term inflation model.

We also proposed a D-term inflation model with a shift symmetry in the imaginary direction of the inflaton superfield and a small linear term in the Kähler potential. We consider the case that the mass of gravitino is $O(10^{2\text{–}3})$ TeV and the mass of the LSP is $O(10^{2\text{–}3})$ GeV, which is naturally realized in anomaly-mediated SUSY-breaking models. In this model, the inflaton decays mainly into gravitinos through the supergravity effects [17–23], and the subsequent decay of gravitinos is a source of nonthermal production of DM. Together with the estimation of baryon asymmetry generated from the Affleck-Dine mechanism, the resulting DM density gives

[8]Note that the inflaton also decays into the MSSM fields with a branching $\sim 10^{-(1\text{–}2)}$, so that there exists significant thermal plasma after the decay of the inflaton. Thus, the sphaleron effect proceeds fast enough to convert the $B-L$ asymmetry to the baryon asymmetry even if the dominant component of the Universe is gravitinos at that time.

[9]Another scenario for cogenesis of baryon and DM has been proposed in Ref. [86], where they introduce an additional heavy field to generate both of them. The decay of Q-balls can also be a solution of the coincidence [35,40–51].

an $\mathcal{O}(1)$ baryon-to-DM ratio. This is because both of them are related to the energy scale of inflation. The amount of baryon asymmetry is proportional to the reheating temperature of the Universe and inversely proportional to the Hubble parameter during inflation. That of DM is proportional to the reheating temperature and inversely proportional to the mass of the inflaton. Since the Hubble parameter and the mass are related to each other, the resulting baryon and DM density is naturally of order unity. We predict that the LSP mass is 2 orders of magnitude larger than the proton mass, which comes from the fact that the GUT scale is 2 orders of magnitude less than the Planck scale. When the LSP is mostly wino or Higgsino, it would be detected by future indirect detection experiments of DM.

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APPENDIX: ANOTHER SOLUTION FOR THE BARYON-DM COINCIDENCE PROBLEM

In this appendix, we consider the case in which the mass of the Polonyi field is less than the Hubble parameter H_I and it obtains a nonzero VEV during inflation. Even if the origin of its potential is a symmetry-enhanced point, the Polonyi field cannot be stabilized at the origin because of the absence of the Hubble-induced mass during *D*-term inflation. Although one might wonder if this is an obstacle known as the Polonyi problem, we show that its decay can explain the amount of DM once we allow a 10% fine-tuning for the initial VEV of the Polonyi field.

During *D*-term inflation, the Polonyi field obtains a VEV denoted as z_i, which might be as large as the Planck scale. Just after inflation ends, the Polonyi field starts to oscillate around the origin due to the positive Hubble-induced mass term as the AD field considered in Sec. III. Then, its number density decreases with time as a^{-3} due to the expansion of the Universe. This means that its energy density decreases as $a^{-9/2}$ because its Hubble-induced mass is of the order of the Hubble parameter, which decreases as $a^{-3/2}$. After the Hubble parameter decreases down to its low-energy mass of the Polonyi field [that is, after the time of $H(t) \simeq m_z$], the Hubble-induced mass term can be neglected. Then its mass and VEV are given by Eqs. (38) and (39), respectively. Note that the minimum of the potential is usually much smaller than the amplitude of the Polonyi field at that time because its amplitude decreases only as $a^{-3/4}$ until $H(t) \simeq m_z$:

$$z|_{H(t)=m_z} \simeq \left(\frac{m_z}{H_I}\right)^{1/2} z_i \gg \langle z \rangle_0. \tag{A1}$$

This means that its number density is not affected and continues to decrease with time as a^{-3}.

The Polonyi field decays mainly into gravitinos with a rate of

$$\Gamma_z(z \to 2\psi_{3/2}) \simeq \frac{1}{96\pi} \frac{m_z^5}{m_{3/2}^2 M_{\rm Pl}^2}. \tag{A2}$$

Since we consider a relatively low reheating temperature to realize the Affleck-Dine baryogenesis, we can neglect the thermal production of gravitinos. Thus, the gravitino abundance is determined by the number density of the Polonyi field. We require that the mass of gravitino be of order $10^{2\text{–}3}$ TeV so that it decays into the MSSM particles before the BBN epoch. Otherwise, the decay particles interact with the light elements and spoil the success of the BBN [78–82]. The decay of gravitinos is a source of nonthermal production of LSP DM. We assume that the thermal relic of the LSP is much smaller than the observed DM abundance. This can be achieved for the case of wino-like LSP with a mass much less than 3 TeV [89,90] or Higgsino-like LSP with a mass much less than 1 TeV. The wino-like LSP is well motivated in models of anomaly-mediated SUSY breaking, where the gravitino mass is naturally as large as $O(100)$ TeV, as we required.

In summary, the LSP DM is nonthermally produced from the decay of the gravitino, which is generated from the decay of the Polonyi field. The Polonyi field is generated by its coherent oscillation just after the end of inflation. We thus obtain the following DM abundance:

$$\begin{aligned} Y_{\rm DM} &\equiv \frac{n_{\rm LSP}}{s} \\ &\simeq \left.\frac{n_{3/2}}{s}\right|_{\psi_{3/2}\,\rm decay} \\ &\simeq \left.\frac{3T_{\rm RH} n_{3/2}}{4\rho_I}\right|_{\rm RH} \\ &\simeq \left.\frac{3T_{\rm RH} n_z}{2\rho_I}\right|_{H=\Gamma_z} \\ &\simeq \left.\frac{3T_{\rm RH} n_z}{2\rho_I}\right|_{H=H_I} \\ &\simeq \frac{T_{\rm RH} z_i^2}{2H_I M_{\rm Pl}^2}, \end{aligned} \tag{A3}$$

where we use $n_z \simeq H_I z_i^2$ at $H = H_I$ in the last line.

Together with the estimation of baryon asymmetry generated by the Affleck-Dine mechanism calculated in

Sec. III B, the above result implies that the baryon-to-DM ratio is given by the following simple relation:

$$\frac{\Omega_b}{\Omega_{\rm DM}} \simeq \frac{4}{23}\epsilon b \frac{m_p}{m_{\rm LSP}}\left(\frac{\phi_i}{z_i}\right)^2. \tag{A4}$$

Thus, we can explain the observed baryon-to-DM ratio when $z_i/\phi_i \sim 0.1$. One might expect that the natural values of their initial VEVs are of the order of the Planck scale. In that case, the result requires a 10% fine-tuning for the value of z_i.

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