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Baryon asymmetry in a heavy moduli scenario

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In some models of supersymmetry breaking, modulus fields are heavy enough to decay before big bang nucleosynthesis. But the large entropy produced via moduli decay significantly dilutes the preexisting baryon asymmetry of the Universe. We study whether the Affleck-Dine mechanism can provide enough baryon asymmetry which survives the dilution, and find several situations in which a desirable amount of baryon number remains after the dilution. The possibility of nonthermal dark matter is also discussed. This provides the realistic cosmological scenario with heavy moduli.

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I. INTRODUCTION

Recent cosmological observations have revealed that the ordinary matter contributes to the only small fraction of the energy density of the Universe, $\Omega_b h^2 \sim 0.022$ in terms of the density parameter, and the remainder comes from "dark" components, dark matter and dark energy, whose contributions are represented by $\Omega_m h^2 \sim 0.13$ and $\Omega_\Lambda \sim 0.72$, respectively [1]. The existence of these dark components indicates physics beyond the standard model such as supersymmetry (SUSY) [2], but on the other hand, the nonstandard physics is also considered to be imprinted in the ordinary matter (baryon) component. From the measurement of cosmic microwave background (CMB) anisotropy and light element abundances predicted by big bang nucleosynthesis (BBN) [3], it is known that the baryon density of the Universe is almost made only of baryons, and antibaryons do not exist. But such an observed amount of baryon asymmetry cannot be generated within the framework of the standard model. Thus, if there is an underlying physics beyond the standard model, the baryon density of the Universe $\Omega_b h^2 \sim 0.022$ should also be explained by violation of the baryon number and CP built in the new physics, as well as the dark matter. In supersymmetric theory, which is one of the best motivated physics beyond the standard model, there is an interesting mechanism to create baryon asymmetry. In SUSY there exist many scalar fields as superpartner of the standard model fermion which carry the baryon or lepton number. Along some directions of the configuration of these scalar fields, the scalar potential is flat. Thus scalar fields corresponding to these flat directions can develop to large field value during inflation, and subsequent evolution of the scalar fields naturally leads to baryon asymmetry. This is called Affleck-Dine mechanism [4], which we will focus on in this paper.

On the other hand, global supersymmetry is naturally extended to the local supersymmetry, which inevitably includes gravity, that is supergravity. In supergravity there appear long-lived massive particles whose lifetimes are typically longer than 1 sec and they decay after BBN starts. One is the gravitino, the superpartner of the graviton. Gravitinos are produced in high-temperature plasma via scatterings of the particles in thermal bath and their subsequent decay may greatly affect the standard cosmology [5,6], or may overclose the Universe if they are stable [7]. Another is the Polonyi field, which is a singlet scalar field introduced in order to give the SUSY breaking masses to the superparticles, especially the mass of gauginos, in many models of SUSY breaking. Generally the Polonyi field has the large field value during inflation, and it begins to oscillate coherently with initial amplitude of order reduced Planck scale M_P when the Hubble parameter H becomes equal to the gravitino mass $m_{3/2}$. It has extremely large energy density and its decay after BBN has catastrophic effects on the standard cosmology [8]. Furthermore, supergravity may be the low-energy effective theory of string theory, which is defined in 10-dimensional spacetime. In compactification of such extra dimensions, there appear light scalar fields called moduli. Generally moduli have the mass of order $m_{3/2}$ through nonperturbative dynamics associated with SUSY breaking. The dynamics of moduli fields and cosmological difficulty they cause are similar to those of the Polonyi field, and we call them the moduli problem [9].

There are some suggestions to solve the moduli problem. One possible way is to use late-time inflation and the subsequent entropy production in order to dilute the moduli abundance to a cosmologically safe value. Such late-time inflation is realized by thermal inflation models and concrete examples for them are found in Refs. [10–12]. The other way is to make moduli heavy enough to decay well before BBN. The modulus mass larger than about 100 TeV is safe and such a large mass is naturally realized in anomaly-mediated SUSY breaking models, where other SUSY particles obtain the mass of order $\sim m_{3/2}/(8\pi^2) \sim 1$ TeV [13].

But the scenario is not complete. In the late-time entropy production scenario, the preexisting baryon asymmetry is also diluted. The reheating temperature after thermal inflation is typically less than a few GeV, and hence almost all baryogenesis mechanisms which rely on high-energy physics do not work. The variant type of the Affleck-Dine mechanism after thermal inflation may work [14] and to the best of our knowledge it is only a possibility to create enough baryon asymmetry in the presence of thermal inflation. In the heavy moduli scenario, significant entropy released by the decay of moduli also dilute the preexisting baryon asymmetry. In previous works [15,16], it was assumed that the Affleck-Dine mechanism can create enough baryon asymmetry which survives the dilution from moduli decay. However, in fact, for large $m_{3/2}$ the ordinary Affleck-Dine mechanism does not work due to the nontrivial potential minima of the Affleck-Dine field [17,18].

In this paper we study whether the sufficient baryon asymmetry is created in the heavy moduli scenario, such as anomaly-mediated SUSY breaking models or mixed modulus anomaly-mediation (or mirage-mediation) models [19]. Mirage-mediation models are based on the concrete model of Kachru-Kallosh-Linde-Trivedi (KKLT) flux compactification [20] and the lightest modulus mass is predicted as $m_\chi \sim (4\pi^2)m_{3/2}$. Thus in the following we consider the typical situations where the modulus mass m_χ is of the order $m_{3/2}$ and $(4\pi^2)m_{3/2}$ as reference values, although we do not specify the concrete SUSY breaking models. The only assumption on which our analysis is based is the hierarchical relation between the gravitino mass and other SUSY particle masses. In both scenarios (anomaly- or mirage-mediation models) other SUSY particles have mass of order $\sim m_{3/2}/(8\pi^2)$ and hence the gravitino is considered to be as heavy as 100 TeV. As we will explain later, the Affleck-Dine mechanism for such large $m_{3/2}$ is highly nontrivial. One possible way to incorporate the Affleck-Dine mechanism in the anomaly- or mirage-mediation model is to introduce the gauged $U(1)_{B-L}$ symmetry [21]. But we pointed out that the Affleck-Dine mechanism with large $m_{3/2}$ works without any additional assumption if the reheating temperature is relatively high [22]. We examine these two scenarios and find a possible parameter region in which the desired amount of baryon asymmetry is generated.

This paper is organized as follows. In Sec. II, we briefly review the cosmological dynamics of moduli fields. In Sec. III, Affleck-Dine baryogenesis in a high reheating temperature scenario is discussed. In Sec. IV, Affleck-Dine baryogenesis in the gauged $U(1)_{B-L}$ model is discussed. Somewhat similar to this case, but the case without superpotential and Q-ball domination is discussed in Sec. V. The decay of moduli may cause another difficulty, especially lightest supersymmetric particle (LSP) overproduction from the decay of moduli-induced gravitinos. We give possible solutions to this problem in Sec. VI. We conclude in Sec. VII.

II. DYNAMICS OF MODULUS FIELDS

In general, a modulus field has the mass of order $m_{3/2}$ in the presence of SUSY breaking. In the early Universe, the inflaton dominates the energy density and this vacuum energy also breaks supersymmetry. As a result the modulus field obtains Hubble-induced SUSY breaking mass of order H [23]. Thus, during inflation the modulus has very large mass and sits at the origin. However, this high-energy minimum does not need to coincide with the low-energy true minimum of the potential. In general, these two minima are expected to be separated by the Planck scale and hence when H becomes smaller than the modulus mass m_χ, the modulus field begins to oscillate with the initial amplitude $\chi_0 \sim M_P$. This coherent oscillation of the modulus field has the energy density $\sim m_\chi^2 M_P^2$ initially, and the total energy density of the Universe is given by $\sim 3m_\chi^2 M_P^2$. We can see that the modulus has inevitably large energy density comparable to the total energy density and dominates the Universe as soon as the inflaton decays (in the case of $m_\chi > \Gamma_I$, where Γ_I denotes the decay rate of the inflaton) or the modulus field begins to oscillate (in the case of $\Gamma_I > m_\chi$).

In the case of $m_\chi > \Gamma_I$ where the inflaton decays after the moduli start to oscillate, the moduli-to-entropy ratio is given by

$$\frac{\rho_\chi}{s} = \frac{m_\chi^2 \chi(T_R)^2}{4\rho(T_R)/3T_R} = \frac{1}{4}T_R\left(\frac{\chi_0}{M_P}\right)^2 \sim 2.5\times 10^4\ \text{GeV}\left(\frac{T_R}{10^5\ \text{GeV}}\right)\left(\frac{\chi_0}{M_P}\right)^2, \tag{1}$$

where T_R denotes the reheating temperature from the inflaton decay. On the other hand, in the case of $m_\chi < \Gamma_I$, it is given by

$$\frac{\rho_\chi}{s} = \frac{1}{4}\left(\frac{90}{\pi^2 g_*}\right)^{1/4} m_\chi^{1/2} M_P^{1/2}\left(\frac{\chi_0}{M_P}\right)^2 \sim 1.2\times 10^{11}\ \text{GeV}\left(\frac{90}{\pi^2 g_*}\right)^{1/4}\left(\frac{m_\chi}{100\ \text{TeV}}\right)^{1/2}\left(\frac{\chi_0}{M_P}\right)^2. \tag{2}$$

In both cases the moduli abundance largely exceeds the critical density of the Universe, $\rho_c/s_0 \sim 3.6\times 10^{-9}h^2$ GeV. If we parametrize the decay width of the modulus as

$$\Gamma_\chi = \frac{c}{4\pi}\frac{m_\chi^3}{M_P^2}, \tag{3}$$

the decay temperature of moduli T_χ is given by

$$T_\chi \sim 5.5\ \text{MeV}\sqrt{c}\left(\frac{m_\chi}{100\ \text{TeV}}\right)^{3/2}. \tag{4}$$

Hence, the decay temperature of modulus typically takes a value from a few MeV to a few GeV for 100 TeV $\lesssim m_\chi \lesssim (4\pi^2)100$ TeV. The late decay of moduli with such large abundance has significant effects on BBN [24], CMB [25], and diffuse X(γ)-ray background [26], which results in strong disagreement with observations. But for the modulus mass larger than about 100 TeV, the moduli decay before BBN and do not spoil the success of standard BBN. The Universe is finally reheated by the moduli with reheating temperature T_χ.[1]

III. AFFLECK-DINE BARYOGENESIS WITH EARLY OSCILLATION

A. The model

In the minimal supersymmetric standard model (MSSM), there exist some configurations of the scalar fields (flat directions) along which scalar potential vanishes in the supersymmetric limit within the renormalizable terms [29]. A flat direction is parametrized by the single complex scalar ϕ, which we call the Affleck-Dine field. The Affleck-Dine field feels potentials from nonrenormalizable superpotentials represented as

$$W_{\rm NR} = \frac{\phi^n}{nM^{n-3}}, \qquad (5)$$

where M is the effective cutoff scale and $n \geq 4$. Including SUSY breaking effects, the potential for the Affleck-Dine field is written as

$$V_S(\phi) = m_\phi^2|\phi|^2 + \left(a_m m_{3/2}\frac{\phi^n}{nM^{n-3}} + {\rm H.c.}\right) + \frac{|\phi|^{2(n-1)}}{M^{2(n-3)}}, \qquad (6)$$

where a_m is $O(1)$ numerical coefficient. There are other sources for the scalar potentials. As explained in Sec. II, the scalar fields obtain Hubble-induced SUSY breaking terms such as

$$V_H(\phi) = -c_H H^2|\phi|^2 + \left(a_H H\frac{\phi^n}{nM^{n-3}} + {\rm H.c.}\right), \qquad (7)$$

where c_H and a_H are $O(1)$ coefficients.[2] Here we assume $c_H > 0$. Furthermore, in the high-temperature environment of the early Universe, thermal corrections to the scalar potential also arise. These are [32,33]

$$V_T(\phi) = \sum_{f_k|\phi|<T} c_k f_k^2 T^2|\phi|^2 + a\alpha(T)^2 T^4 \log\left(\frac{|\phi|^2}{T^2}\right), \qquad (8)$$

where c_k is a constant of order unity, f_k denotes gauge or Yukawa couplings relevant for the Affleck-Dine field, and a is a constant of order unity assumed to be positive, which is determined by the two-loop finite-temperature effective potential for the Affleck-Dine field. Then the total scalar potential for the Affleck-Dine field is the sum of them,

$$V(\phi) = V_S(\phi) + V_H(\phi) + V_T(\phi). \qquad (9)$$

Let us summarize the dynamics of the Affleck-Dine field. First, it is trapped by the minimum determined by the balance of the negative Hubble-induced mass term and nonrenormalizable term,

$$|\phi| \simeq (HM^{n-3})^{1/(n-2)}, \qquad (10)$$

and tracks this minimum as H becomes small. The important fact is that without the finite-temperature effect, the scalar potential has the global minimum at

$$|\phi|_{\rm min} \sim \left(\frac{|a_m|}{n-1}m_{3/2}M^{n-3}\right)^{1/(n-2)} \qquad (11)$$

if $m_{3/2}$ is much greater than m_ϕ, which is the situation we are interested in. Thus the Affleck-Dine field is eventually trapped by this minimum and leads to a charge or color breaking vacuum, which is a disaster. But finite-temperature effects can save the situation. Including finite-temperature effects, when H becomes equal to $H_{\rm os}$ determined by

$$H_{\rm os}^2 \sim m_\phi^2 + \sum_{f_k|\phi|<T} c_k f_k^2 T^2 + a\alpha(T)^2\frac{T^4}{|\phi|^2}, \qquad (12)$$

the Affleck-Dine field begins to oscillate around its minimum of the potential. The important fact is that if the thermal log term dominates the potential and oscillation begins by this term, the Affleck-Dine field will be taken to the origin without being trapped by the global minimum [22]. Through the process of field evolution, the Affleck-Dine field receives angular kick from A terms and results in elliptical motion around the origin. Hence the baryon number is generated and conserved in the comoving volume. In fact, as we will see, for high reheating temperature from the inflaton and high field value, the oscillation starts when the thermal logarithmic term dominates the potential and the Affleck-Dine mechanism works well. Note that although the resultant vacuum is metastable, the lifetime of the false vacuum is much longer than the age of the Universe [18].

B. Baryon asymmetry

Next we estimate the baryon asymmetry in the presence of a heavy modulus field. In the case of early oscillation due to thermal logarithmic potential, $H_{\rm os}$ is given by

$$H_{\rm os} = \alpha T_R\left(\frac{M_P}{M}\right)^{1/2} \quad \text{for } n = 4, \qquad (13)$$

$$H_{\rm os} = (\alpha^2 T_R^2 M_P M^{-3/2})^{2/3} \quad \text{for } n = 6. \qquad (14)$$

Hereafter we consider the $n = 4$ and $n = 6$ case only, because flat directions with $n \geq 7$ are lifted by the super-

[1] If the modulus field does not have the Hubble mass and obtain unsuppressed quantum fluctuation during inflation, it can be the interesting candidate of the curvaton [27,28]. But we do not go into the details of this issue.

[2] In some inflation models such as D-term inflation [30], the Hubble-induced term does not arise during inflation [31].

potential of the form $\psi\phi^{n-1}/M^{n-3}$, where ψ represents the scalar field other than the Affleck-Dine field and cannot generate baryon asymmetry via the Affleck-Dine mechanism [29]. We need the following constraints for this scenario to work. One is the condition that early oscillation occurs ($H_{\rm os} > m_\phi$), this leads to

$$T_R \gtrsim \frac{m_\phi}{\alpha}\left(\frac{M}{M_P}\right)^{1/2} \sim 1\times 10^3\ \mathrm{GeV}\left(\frac{0.1}{\alpha}\right)\left(\frac{m_\phi}{100\ \mathrm{GeV}}\right)\left(\frac{M}{M_P}\right)^{1/2} \tag{15}$$

for $n = 4$, and

$$T_R \gtrsim \frac{1}{\alpha}\left(\frac{m_\phi^3 M^3}{M_P^2}\right)^{1/4} \sim 3\times 10^5\ \mathrm{GeV}\left(\frac{0.1}{\alpha}\right)\left(\frac{m_\phi}{100\ \mathrm{GeV}}\right)^{3/4}\left(\frac{M}{10^{15}\ \mathrm{GeV}}\right)^{3/4} \tag{16}$$

for $n = 6$. The other is the condition that thermal correction hides the valley of the potential around the true charge-breaking minimum and this thermal logarithmic potential leads the Affleck-Dine field to the origin. This condition is written in the form $\alpha^2 T_{\rm os}^4 \gtrsim |V(|\phi|_{\rm min})|$, explicitly,

$$T_R \gtrsim \alpha^{-1} m_{3/2}\left(\frac{M_P}{M}\right)^{1/2} \sim 1\times 10^6\ \mathrm{GeV}\left(\frac{0.1}{\alpha}\right)\left(\frac{m_{3/2}}{100\ \mathrm{TeV}}\right)\left(\frac{M}{M_P}\right)^{1/2} \tag{17}$$

for $n = 4$ and

$$T_R \gtrsim \alpha^{-1} M^{3/4} M_P^{-1/2} m_{3/2}^{3/4} \sim 8\times 10^5\ \mathrm{GeV}\left(\frac{0.1}{\alpha}\right)\left(\frac{m_{3/2}}{10\ \mathrm{TeV}}\right)^{3/4}\left(\frac{M}{10^{15}\ \mathrm{GeV}}\right)^{3/4} \tag{18}$$

for $n = 6$. We can see that in general the latter condition is severer whenever $m_{3/2} > m_\phi$, which is always satisfied in anomaly-mediated SUSY breaking.

Now let us estimate the baryon asymmetry in the presence of modulus fields. It is convenient to express the baryon-to-entropy ratio as

$$\frac{n_B}{s} = \frac{n_B}{\rho_\chi}\frac{\rho_\chi(T_\chi)}{s(T_\chi)} = \frac{n_B}{\rho_\chi}\frac{3T_\chi}{4}. \tag{19}$$

The ratio n_B/ρ_χ is fixed when both the Affleck-Dine field and modulus field begin to oscillate. When $H_{\rm os} > m_\chi$, this ratio is fixed at the onset of the oscillation of the moduli, $H = m_\chi$. On the other hand, when $H_{\rm os} < m_\chi$, the ratio is fixed at the beginning of the Affleck-Dine field oscillation, $H = H_{\rm os}$. From Eqs. (14) and (17) or (18),

$$H_{\rm os} \gtrsim m_{3/2} \tag{20}$$

must always hold. Thus if we assume $m_\chi \sim m_{3/2}$ we can safely focus on the case $H_{\rm os} > m_\chi$. But in some models based on string theory, $m_\chi \gg m_{3/2}$ might be possible. In the mirage-mediation model, the modulus mass is predicted as $m_\chi \sim 4\pi^2 m_{3/2}$ [19]. Although in such a model $H_{\rm os} < m_\chi$ is still possible, we mainly focus on the case $H_{\rm os} > m_\chi$ and briefly discuss the modification in the case $H_{\rm os} < m_\chi$. The condition $H_{\rm os} > m_\chi$ is rewritten as

$$T_R \gtrsim \alpha^{-1} m_\chi\left(\frac{M_P}{M}\right)^{1/2} \sim 1\times 10^6\ \mathrm{GeV}\left(\frac{0.1}{\alpha}\right)\left(\frac{m_\chi}{100\ \mathrm{TeV}}\right)\left(\frac{M}{M_P}\right)^{1/2} \tag{21}$$

for $n = 4$ and

$$T_R \gtrsim \alpha^{-1} m_\chi^{3/4} M^{3/4} M_P^{-1/2} \sim 2\times 10^9\ \mathrm{GeV}\left(\frac{0.1}{\alpha}\right)\left(\frac{m_\chi}{100\ \mathrm{TeV}}\right)^{3/4}\left(\frac{M}{M_P}\right)^{3/4} \tag{22}$$

for $n = 6$. High reheating temperature from the inflaton is not a problem as far as the moduli decay well before BBN and nonthermal LSPs associated with modulus decay do not overclose the Universe. In fact it is possible that nonthermal LSPs from the decay of moduli account for the present matter density of the Universe (see Sec. VI).

1. $H_{\rm os} > m_\chi$

In this case the baryon-to-moduli ratio n_B/ρ_χ is fixed at $H = m_\chi$ where the modulus field begins to oscillate with amplitude $\chi_0 \sim M_P$,

$$\frac{n_B}{\rho_\chi} = \frac{n_B(t_{\rm os})}{m_\chi^2\chi_0^2}\left(\frac{a(t_{\rm os})}{a(t_{\rm mod})}\right)^3, \tag{23}$$

where $t_{\rm mod} \simeq m_\chi^{-1}$. In order to get the correct estimation, we must specify the decay epoch of the inflaton, whose decay rate is denoted as Γ_I. Thus depending on Γ_I, three scenarios are available: (a) $\Gamma_I > H_{\rm os} > m_\chi$, (b) $H_{\rm os} > \Gamma_I > m_\chi$, (c) $H_{\rm os} > m_\chi > \Gamma_I$. Note that in case (a), at the beginning of oscillation of the Affleck-Dine field the Universe already enters the radiation dominated era and estimation of baryon number is somewhat different from the other two cases. Before the estimation, we see the conditions when the case (a), (b), and (c) are realized. The condition that $m_\chi > \Gamma_I$ can be written as

$$T_R \lesssim 2\times 10^{11}\ \mathrm{GeV}\left(\frac{m_\chi}{100\ \mathrm{TeV}}\right)^{1/2}. \tag{24}$$

On the other hand, the condition $H_{\rm os} > \Gamma_I$ can be rewritten as follows,

$$T_R \lesssim 5\times 10^{16}\ \mathrm{GeV}\left(\frac{\alpha}{0.1}\right)\left(\frac{M_P}{M}\right)^{1/2} \quad \text{for } n = 4, \qquad T_R \lesssim 2\times 10^{15}\ \mathrm{GeV}\left(\frac{\alpha}{0.1}\right)^2\left(\frac{M_P}{M}\right)^{3/2} \quad \text{for } n = 6 \tag{25}$$

which is satisfied for a natural range of parameters. In other words, unless the reheating temperature is unnaturally high, case (a) is not realized. The conditions (21) (or (22)), (24) and (25) determine which of the following scenario is realized.

In the case (a), early oscillation begins in the radiation dominated regime and the baryon-to-moduli ratio (23) is written as

$$\frac{n_B}{\rho_\chi} = \frac{\delta_e m_{3/2}|\phi_{\rm os}|^2}{m_\chi^2 \chi_0^2}\left(\frac{m_\chi}{H_{\rm os}}\right)^{3/2}, \tag{26}$$

where $\delta_e(\sim O(1))$ denotes the effective CP phase. Note that as far as the initial amplitude of the Affleck-Dine field is smaller than M_P, it never dominates the Universe at the instant of oscillation. We can estimate $\phi_{\rm os}$ and $H_{\rm os}$ as

$$|\phi|_{\rm os} \sim \gamma_*^{-(1/2)} \alpha M_P \tag{27}$$

and

$$H_{\rm os} \sim \frac{\alpha^2 M_P^2}{\gamma_* M} \quad \text{for } n = 4, \qquad H_{\rm os} \sim \frac{\alpha^4 M_P^4}{\gamma_*^2 M^3} \quad \text{for } n = 6, \tag{28}$$

where $\gamma_* = (\pi^2 g_*(T_R)/90) \sim 25$. Substituting these values and using Eq. (4), we finally obtain the baryon-to-entropy ratio after decay of the modulus field as

$$\frac{n_B}{s} = \frac{0.2\delta_e\sqrt{c\gamma_*}}{\alpha}\frac{m_{3/2}m_\chi}{M_P^2}\left(\frac{M}{M_P}\right)^{3/2}\left(\frac{M_P}{\chi_0}\right)^2 \sim 7\times 10^{-27}\delta_e\sqrt{c}\left(\frac{0.1}{\alpha}\right)\left(\frac{m_{3/2}}{100\ \text{TeV}}\right)\left(\frac{m_\chi}{100\ \text{TeV}}\right) \times \left(\frac{M}{M_P}\right)^{3/2}\left(\frac{M_P}{\chi_0}\right)^2 \tag{29}$$

in the case of $n = 4$ flat direction. Clearly this is too small and it is impossible that we obtain a proper amount of baryon asymmetry. For $n = 6$, we obtain

$$\frac{n_B}{s} = \frac{0.2\delta_e\sqrt{c}\gamma_*^2}{\alpha^4}\frac{m_{3/2}m_\chi}{M_P^2}\left(\frac{M}{M_P}\right)^{9/2}\left(\frac{M_P}{\chi_0}\right)^2 \sim 2\times 10^{-21}\sqrt{c}\delta_e\left(\frac{0.1}{\alpha}\right)^4\left(\frac{m_{3/2}}{100\ \text{TeV}}\right)\left(\frac{m_\chi}{100\ \text{TeV}}\right) \times \left(\frac{M}{M_P}\right)^{9/2}\left(\frac{M_P}{\chi_0}\right)^2. \tag{30}$$

It also seems too small, but dependence of the cutoff scale M is rather large, so that if we assume $M \sim 100 M_P$ the desired amount of baryon asymmetry can be obtained. It may seem peculiar that the cutoff scale M is bigger than the Planck scale, but our definition of M includes some coupling constant, e.g., even if the physical cutoff scale is M_P, the effective cutoff scale can be $\sim 100 M_P$ if the relevant coupling constant is 10^{-2}.

In the case (b), the modulus oscillation begins in the radiation dominated era. The baryon-to-moduli ratio (23) is expressed as

$$\frac{n_B}{\rho_\chi} = \frac{\delta_e m_{3/2}|\phi_{\rm os}|^2}{m_\chi^2\chi_0^2}\left(\frac{\Gamma_I}{H_{\rm os}}\right)^2\left(\frac{m_\chi}{\Gamma_I}\right)^{3/2}. \tag{31}$$

A straightforward calculation yields

$$\frac{n_B}{s} = \frac{0.5\delta_e\sqrt{c}}{\alpha}\frac{m_{3/2}m_\chi}{M_P^2}\left(\frac{M}{M_P}\right)^{3/2}\left(\frac{M_P}{\chi_0}\right)^2 \sim 8\times 10^{-27}\delta_e\sqrt{c}\left(\frac{0.1}{\alpha}\right)\left(\frac{m_{3/2}}{100\ \text{TeV}}\right)\left(\frac{m_\chi}{100\ \text{TeV}}\right)\left(\frac{M}{M_P}\right)^{3/2} \times \left(\frac{M_P}{\chi_0}\right)^2 \tag{32}$$

for the $n = 4$ case. Obviously, this is too small. On the other hand, for the $n = 6$ case we obtain

$$\frac{n_B}{s} = \frac{0.5\delta_e\sqrt{c}}{\alpha^2}\frac{m_{3/2}m_\chi}{T_R M_P}\left(\frac{M}{M_P}\right)^3\left(\frac{M_P}{\chi_0}\right)^2 \sim 2\times 10^{-16}\delta_e\sqrt{c}\left(\frac{0.1}{\alpha}\right)^2\left(\frac{m_{3/2}}{100\ \text{TeV}}\right)\left(\frac{m_\chi}{100\ \text{TeV}}\right) \times \left(\frac{10^9\ \text{GeV}}{T_R}\right)\left(\frac{M}{M_P}\right)^3\left(\frac{M_P}{\chi_0}\right)^2. \tag{33}$$

It seems possible that a proper amount of baryon asymmetry is generated after choosing the cutoff scale appropriately. But T_R is constrained from the condition $H_{\rm os} > m_\chi$ [Eq. (22)]. Substituting Eq. (22) into the above equation, we obtain the upper bound on n_B/s,

$$\frac{n_B}{s} \lesssim \frac{0.5\delta_e\sqrt{c}}{\alpha}\frac{m_{3/2}^{1/4}m_\chi}{M_P^{5/4}}\left(\frac{M}{M_P}\right)^{9/4}\left(\frac{M_P}{\chi_0}\right)^2 \sim 8\times 10^{-17}\delta_e\sqrt{c}\left(\frac{0.1}{\alpha}\right)^2\left(\frac{m_{3/2}}{100\ \text{TeV}}\right)^{1/4}\left(\frac{m_\chi}{100\ \text{TeV}}\right) \times \left(\frac{M}{M_P}\right)^{9/4}\left(\frac{M_P}{\chi_0}\right)^2. \tag{34}$$

Thus we need $M \gtrsim 200 M_P$ to obtain enough baryon asymmetry. If this is the case, T_R must also be as high as 10^{12} GeV. This also satisfies the constraint $\Gamma_I > m_\chi$.

In the case (c), the modulus starts to oscillate in the inflaton-dominated regime and then the inflaton decays resulting in the brief radiation dominated era followed by the moduli dominated universe. The baryon-to-moduli ratio (23) in this case is expressed as

$$\frac{n_B}{\rho_\chi} = \frac{\delta_e m_{3/2}|\phi_{\rm os}|^2}{m_\chi^2\chi_0^2}\left(\frac{m_\chi}{H_{\rm os}}\right)^2. \tag{35}$$

For the $n = 4$ case, we obtain

$$\frac{n_B}{s} = \frac{0.2\delta_e\sqrt{c}}{\alpha}\frac{m_{3/2}m_\chi^{3/2}}{T_R M_P^{3/2}}\left(\frac{M}{M_P}\right)^{3/2}\left(\frac{M_P}{\chi_0}\right)^2 \sim 2\times 10^{-24}\delta_e\sqrt{c}\left(\frac{0.1}{\alpha}\right)\left(\frac{m_{3/2}}{100\text{ TeV}}\right)\left(\frac{m_\chi}{100\text{ TeV}}\right)^{3/2} \times\left(\frac{10^9\text{ GeV}}{T_R}\right)\left(\frac{M}{M_P}\right)^{3/2}\left(\frac{M_P}{\chi_0}\right)^2, \quad (36)$$

which is extremely small compared with the present baryon density. When we apply to the $n = 6$ flat direction, the baryon-to-entropy ratio is estimated as

$$\frac{n_B}{s} = \frac{0.2\delta_e\sqrt{c}}{\alpha^2}\frac{m_{3/2}m_\chi^{3/2}}{T_R^2 M_P^{1/2}}\left(\frac{M}{M_P}\right)^{3}\left(\frac{M_P}{\chi_0}\right)^2 \sim 4\times 10^{-14}\delta_e\sqrt{c}\left(\frac{0.1}{\alpha}\right)^2\left(\frac{m_{3/2}}{100\text{ TeV}}\right)\left(\frac{m_\chi}{100\text{ TeV}}\right)^{3/2} \times\left(\frac{10^9\text{ GeV}}{T_R}\right)^2\left(\frac{M}{M_P}\right)^{3}\left(\frac{M_P}{\chi_0}\right)^2, \quad (37)$$

which seems successful. However, we need rather high reheating temperature which suppresses the baryon-to-entropy ratio, due to the condition $H_{os} > m_\chi$ [Eq. (22)]. Substituting Eq. (22), the upper limit for baryon asymmetry is obtained,

$$\frac{n_B}{s} \lesssim 0.2\delta_e\sqrt{c}\frac{m_{3/2}}{M_P}\left(\frac{M}{M_P}\right)^{3/2} \sim 8\times 10^{-15}\delta_e\sqrt{c}\left(\frac{0.1}{\alpha}\right)^2\left(\frac{m_{3/2}}{100\text{ TeV}}\right)\left(\frac{M}{M_P}\right)^{3/2}. \quad (38)$$

If $M \sim 100 M_P$ we can obtain the desired baryon asymmetry, and this indicates that the reheating temperature should be higher than $\sim 10^{11}$ GeV. On the other hand, $T_R \lesssim 10^{12}$ GeV is necessary in order to satisfy $\Gamma_I < m_\chi$. Thus $M \sim 100 M_P$ and 10^{11} GeV $\lesssim T_R \lesssim 10^{12}$ GeV are the possible parameter regions (see Fig. 1).

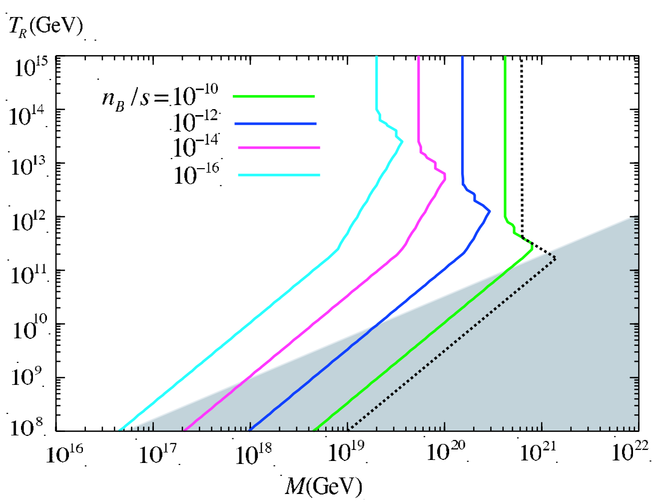

FIG. 1 (color online). The contour plot of n_B/s as a function of M and T_R. We take $m_\chi = m_{3/2} = 100$ TeV. In the shaded region the Affleck-Dine field is trapped into charge-breaking minima and baryogenesis does not work. Also we show by the dotted line $Q \sim 10^{20}$. The left side of the dotted line predicts $Q < 10^{20}$ and Q balls completely evaporate in high-temperature plasma.

2. $H_{os} < m_\chi$

Now let us turn to the case $H_{os} < m_\chi$. As we explained, early oscillation to avoid charge or color breaking minima requires $H_{os} > m_{3/2}$ and hence this particular possibility arises only when modulus mass m_χ is much heavier than $m_{3/2}$.[3] In this case, we can classify the cosmological scenario depending on the inflaton decay rate Γ_I: (d) $\Gamma_I > m_\chi > H_{os}$, (e) $m_\chi > \Gamma_I > H_{os}$, (f) $m_\chi > H_{os} > \Gamma_I$. Note that baryon-to-moduli ratio is fixed once the Affleck-Dine field starts to oscillate, but the resulting formula for n_B/ρ_χ is the same as Eq. (23). Therefore the results of the case (d) and (f) are the same as (a) and (c), respectively. Only the case (e) slightly differs from (b).

In the case (e) the baryon-to-moduli ratio is expressed as

$$\frac{n_B}{\rho_\chi} = \frac{\delta_e m_{3/2}|\phi_{os}|^2}{m_\chi^2\chi_0^2}\left(\frac{m_\chi}{\Gamma_I}\right)^2\left(\frac{\Gamma_I}{H_{os}}\right)^2, \quad (39)$$

which is slightly different from the case (b). Note that we have used the approximation that the moduli dominate the Universe soon after the oscillation. The following calculations are similar, and the result is

$$\frac{n_B}{s} = \frac{0.2\delta_e\sqrt{c}\gamma_*^4}{\alpha^6}\frac{m_{3/2}m_\chi^{3/2}T_R^4}{M_P^{13/2}}\left(\frac{M}{M_P}\right)^{4}\left(\frac{M_P}{\chi_0}\right)^2 \gtrsim 4\times 10^{-30}\delta_e\sqrt{c}\left(\frac{0.1}{\alpha}\right)^2\left(\frac{m_{3/2}}{100\text{ TeV}}\right)\left(\frac{m_\chi}{100\text{ TeV}}\right)^{3/2} \times\left(\frac{M}{M_P}\right)^{2}\left(\frac{M_P}{\chi_0}\right)^2 \quad (40)$$

for the $n = 4$ case, where we have used the constraint $\Gamma_I > H_{os}$ in the second line. Using the same constraint, for the $n = 6$ case we obtain

$$\frac{n_B}{s} = \frac{0.2\delta_e\sqrt{c}\alpha^{18}}{\gamma_*^{12}}\frac{m_{3/2}m_\chi^{3/2}M_P^{19/2}}{T_R^{12}}\left(\frac{M_P}{M}\right)^{12}\left(\frac{M_P}{\chi_0}\right)^2 \lesssim 4\times 10^{-33}\delta_e\sqrt{c}\left(\frac{0.1}{\alpha}\right)^6\left(\frac{m_{3/2}}{100\text{ TeV}}\right)\left(\frac{m_\chi}{100\text{ TeV}}\right)^{3/2} \times\left(\frac{M}{M_P}\right)^{6}\left(\frac{M_P}{\chi_0}\right)^2. \quad (41)$$

Similar to the case for $H_{os} > m_\chi$, for the appropriate

[3]As we explain in Sec. VI, although moduli decay into gravitinos may cause cosmological difficulty, here gravitinos are also heavy enough to decay well before BBN. Furthermore LSPs produced by the decay of moduli effectively annihilate and do not overclose the Universe (or they become dark matter). However, the subsequent decay of nonthermally produced gravitinos may pose a cosmological difficulty. See Sec. VI.

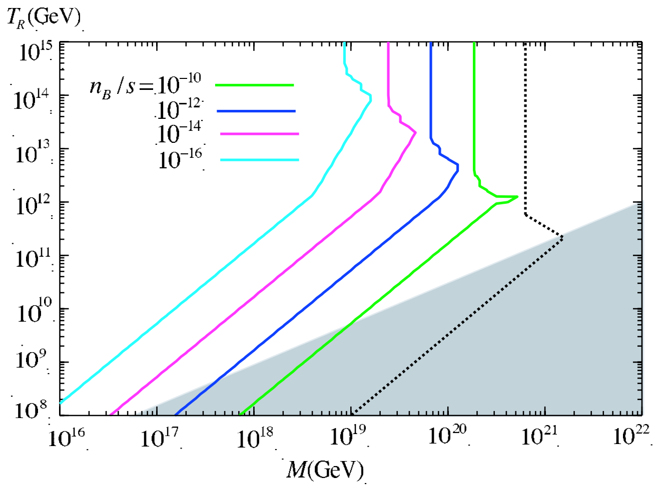

FIG. 2 (color online). Same as Fig. 1, except for $m_\chi = (4\pi^2)m_{3/2} = (4\pi^2)100$ TeV.

choice of the cutoff scale M and the reheating temperature T_R, it seems that we can obtain a proper amount of baryon asymmetry. However, we should recall that the constraint $m_{3/2} < H_{\rm os} < m_\chi$ narrows the allowed parameter range. In fact, the case (e) is not realized in the parameter region we are interested in.

In Figs. 1 and 2 we show the resulting baryon-to-entropy ratio in the (M, T_R) plane in the case of $m_{3/2} = m_\chi = 100$ TeV and $m_\chi = (4\pi^2)m_{3/2} = (4\pi^2)100$ TeV for $n = 6$. The latter case is naturally realized in the mirage-mediation models. We can see that $T_R \gtrsim 10^{11}$ GeV and $M \gtrsim 10^{20}$ GeV are required in the former case. In the latter case where the modulus field is much heavier than the gravitino, the constraint is weaker. Note that in such a heavy moduli scenario gravitinos can be efficiently produced by the decay of moduli, and these nonthermal gravitinos also decay before BBN for $m_{3/2} \sim 100$ TeV. LSPs produced by the decay of those gravitinos may be harmful. We will discuss it in Sec. VI.

C. Q-ball formation

Finally we must consider the effects of Q-ball formation. The fluctuations of the Affleck-Dine field with $U(1)_B$ charge grow and result in lumped condensate, called Q balls [34,35]. The Q-ball formation leads to many nontrivial cosmological consequences, and they highly depend on SUSY breaking models [36,37] (see also [17,38]). As we have seen in the previous subsection, quite a large cutoff scale M is required. One may wonder if this leads to large Q balls and invalidates the applicability of our scenario. However, as we will see, largeness of Q balls is suppressed because of early oscillation. The radius of Q balls is comparable to the Hubble horizon scale at the epoch of Q-ball formation. Thus although the larger cutoff scale M tends to create larger Q balls, but higher reheating temperature T_R, which causes earlier oscillation, tends to make Q balls smaller. Now let us estimate Q.

It is found that for the Q balls which have developed via logarithmic potential, the total charge of the Q ball Q is fitted by the formula [38],

$$Q = \beta\left(\frac{|\phi_{\rm os}|}{T_{\rm os}}\right)^4, \tag{42}$$

where $\beta \sim 6 \times 10^{-4}$. Applying to the early oscillation case for the $n = 6$ flat direction, it is estimated as

$$Q \sim 4 \times 10^{17}\left(\frac{\beta}{6\times 10^{-4}}\right)\left(\frac{10^{11}\ {\rm GeV}}{T_R}\right)^2\left(\frac{M}{100M_P}\right)^3 \tag{43}$$

for $H_{\rm os} > \Gamma_I$, and

$$Q \sim 9 \times 10^{16}\left(\frac{\beta}{6\times 10^{-4}}\right)\left(\frac{0.1}{\alpha}\right)^4\left(\frac{M}{100M_P}\right)^6 \tag{44}$$

for $H_{\rm os} < \Gamma_I$. It is known that evaporation of Q balls in high-temperature plasma can efficiently transfer the charge of Q balls up to $\Delta Q \sim 10^{20}$ almost model independently [39] (see also [22]). Therefore in the most interesting parameter region, Q-balls formed through the Affleck-Dine mechanism can completely evaporate and have no further effects on cosmological evolution of baryon asymmetry.

In Figs. 1 and 2, we show the contour of $Q \sim 10^{20}$ with the black dotted line. It can be seen that in the interesting parameter region where $n_B/s \sim 10^{-10}$ is obtained, only small Q balls are produced and they evaporate in the high-temperature plasma.

IV. AFFLECK-DINE BARYOGENESIS WITH GAUGED $U(1)_{B-L}$

Next we turn to another possibility that Affleck-Dine baryogenesis with large gravitino mass works with an extension of the MSSM to include some additional fields and gauged $U(1)_{B-L}$ symmetry. Because the global $U(1)_{B-L}$ symmetry within MSSM is anomaly-free, it can naturally be extended to local symmetry. But from the viewpoint of baryogenesis, it must be spontaneously broken at some high-energy scale in order to create baryon asymmetry and not to contradict with terrestrial experiments such as proton decay.

A. The model

We briefly explain the model discussed in Ref. [21]. First, we introduce the MSSM singlet fields which have the superpotential as

$$W = \lambda X(S\bar{S} - v^2), \tag{45}$$

where X, S, and $\bar{S}$ have the $U(1)_{B-L}$ charge 0, 2, and -2, respectively, and v denotes the $U(1)_{B-L}$ breaking scale. They induce the scalar potential given by

$$V = |\lambda|^2\{|X|^2(|S|^2 + |\bar{S}|^2) + |S\bar{S} - v^2|^2\} + \frac{g^2}{2}(2|S|^2 - 2|\bar{S}|^2 - q|\phi|^2)^2, \quad (46)$$

where g denotes the $U(1)_{B-L}$ gauge coupling constant and q denotes the $U(1)_{B-L}$ charge of the Affleck-Dine field. The second term comes from the D-term contribution. In the following, we consider flat directions which are lifted by the $n = 6$ nonrenormalizable superpotential in the MSSM, such as udd or LLe direction. In this model, the gauge-invariant superpotential which lifts those flat directions are given by

$$\frac{k_1}{6M^3}\left(\frac{S}{M}\right)(udd)^2, \qquad \frac{k_2}{6M^3}\left(\frac{S}{M}\right)(LLe)^2, \quad (47)$$

where k_1 and k_2 are $O(1)$ coupling constants, and the resulting zero-temperature scalar potential is written as

$$V = m_\phi^2|\phi|^2 - c_{\rm H}H^2|\phi|^2 + \frac{m_{3/2}}{6M^3}\left(\frac{S}{M}\right)(a_m\phi^6 + \text{H.c.}) + \frac{H}{6M^3}\left(\frac{S}{M}\right)(a_H\phi^6 + \text{H.c.}) + \frac{1}{M^6}\left(\frac{S}{M}\right)^2|\phi|^{10} + \frac{1}{36M^8}|\phi|^{12}. \quad (48)$$

Although the whole dynamics is somewhat complicated and we do not give the details here (see [21] for a detail), the point is that by using the additional D-term potential which does not exist in the MSSM, the Affleck-Dine field can be stopped at the $U(1)_{B-L}$ breaking scale v during inflation. If v is smaller than the hill of the potential of the Affleck-Dine field

$$v \lesssim |\phi|_{\rm hill} \sim \left(\frac{m_\phi^2 M^4}{m_{3/2}\langle S\rangle}\right)^{1/4}, \quad (49)$$

the Affleck-Dine mechanism works without trapping into the charge or color breaking global minimum. If we assume $\langle S\rangle \sim v$ and we focus on the $n = 6$ case, this condition is equivalent to

$$v \lesssim \left(\frac{m_\phi^2 M^4}{m_{3/2}}\right)^{1/5} \sim 8\times 10^{14}\ \text{GeV}\left(\frac{100\ \text{TeV}}{m_{3/2}}\right)^{1/5}\left(\frac{m_\phi}{1\ \text{TeV}}\right)^{2/5}\left(\frac{M}{M_P}\right)^{4/5}. \quad (50)$$

If the value v exceeds this bound, the Affleck-Dine baryogenesis cannot work due to the trapping of the Affleck-Dine field in global charge-breaking minima, if thermal effects are neglected.

B. Baryon asymmetry

We saw that in this type of model, the Affleck-Dine field stops at the $U(1)_{B-L}$ breaking scale v until the Hubble parameter becomes of the order m_ϕ and oscillation begins. If v is smaller than the hill of the potential of the Affleck-Dine field, the Affleck-Dine mechanism works. In the case of early oscillation, the result is the same as the usual early oscillation scenario considered in the previous section. Thus we consider only the case of no early oscillation in this subsection. The condition to avoid early oscillation is

$$T_R \lesssim \frac{m_\phi^{1/2}v}{\alpha M_P^{1/2}} \sim 2.1\times 10^8\ \text{GeV}\left(\frac{0.1}{\alpha}\right)\left(\frac{m_\phi}{1\ \text{TeV}}\right)^{2/5}\left(\frac{v}{10^{15}\ \text{GeV}}\right). \quad (51)$$

Thus we can safely set $\Gamma_I < m_\chi$. The baryon number at the instant of oscillation of the Affleck-Dine field is given by

$$n_B(t_{\rm os}) = \frac{4\beta|a_m|}{9}\frac{\delta_e m_{3/2}}{H_{\rm os}M^4}v^7 \quad (52)$$

with $H_{\rm os} \sim m_\phi$. The baryon-to-moduli ratio is once fixed at the epoch of oscillation of the Affleck-Dine field, $t = t_{\rm os}$, and the final reheating comes from the decay of moduli. The result is

$$\frac{n_B}{s} = 0.1\sqrt{c}\delta_e\frac{m_{3/2}m_\chi^{3/2}v^7}{m_\phi^3 M^4 M_P^{5/2}}\left(\frac{M_P}{\chi_0}\right)^2, \quad (53)$$

which depends on the seventh powers of v. Substituting the upper bound on v [Eq. (50)], we obtain an upper bound on the baryon-to-entropy ratio,

$$\frac{n_B}{s} \lesssim 0.1\sqrt{c}\delta_e\frac{m_\chi^{3/2}M^{8/5}}{m_\phi^{1/5}m_{3/2}^{2/5}M_P^{5/2}}\left(\frac{M_P}{\chi_0}\right)^2 \sim 5\times 10^{-13}\sqrt{c}\delta_e\left(\frac{m_\chi}{100\ \text{TeV}}\right)^{3/2}\left(\frac{100\ \text{TeV}}{m_{3/2}}\right)^{2/5} \times\left(\frac{1\ \text{TeV}}{m_\phi}\right)^{1/5}\left(\frac{M}{M_P}\right)^{8/5}\left(\frac{M_P}{\chi_0}\right)^2, \quad (54)$$

which seems successful. However, it is nontrivial whether the Q ball is small enough to evaporate completely. Charge of the Q ball is given by [40]

$$Q \sim \gamma\left(\frac{v}{m_\phi}\right)^2 \times \begin{cases} \epsilon & (\epsilon \gtrsim 0.01) \\ 0.01 & (\epsilon \lesssim 0.01)\end{cases}, \quad (55)$$

where γ is order 10^{-2}–10^{-3} factor which represents the delay of Q-ball formation from the oscillation of Affleck-Dine field and ϵ is called the ellipticity parameter given by

$$\epsilon \sim \delta_e\frac{m_{3/2}v^5}{m_\phi^2 M^4}.$$

Therefore, using the upper bound of v [Eq. (50)], we obtain

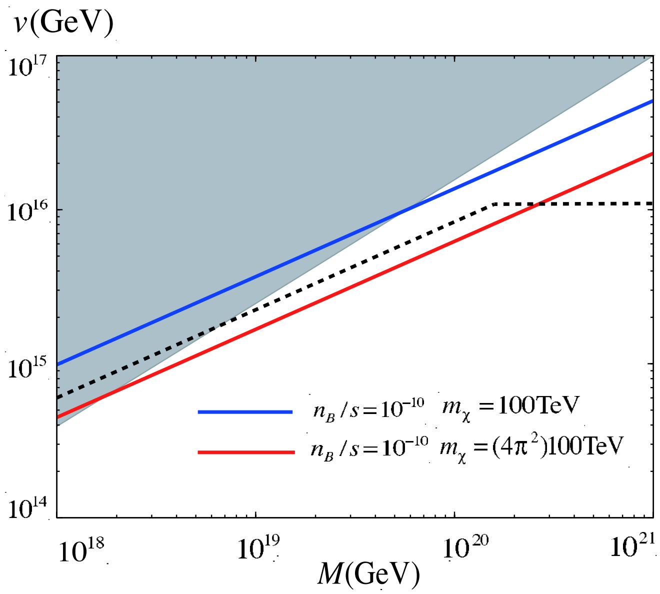

FIG. 3 (color online). The two solid lines show $n_B/s \sim 10^{-10}$ in the gauged $U(1)_{B-L}$ scenario, the upper (blue) line corresponds to $m_\chi = 100$ TeV and the lower (red) line corresponds to $m_\chi = (4\pi^2)100$ TeV. We take $m_{3/2} = 100$ TeV. In the dark shaded region the Affleck-Dine field is trapped into charge-breaking minima and baryogenesis does not work. We also show $Q \sim 10^{22}$ by the dotted line.

$$Q \sim \frac{4\delta_e\gamma}{9}\frac{m_{3/2}v^7}{m_\phi^4 M^4} \lesssim 1\times 10^{21}\delta_e\left(\frac{\gamma}{6\times 10^{-3}}\right)\left(\frac{100\ \mathrm{TeV}}{m_{3/2}}\right)^{2/5}\left(\frac{1\ \mathrm{TeV}}{m_\phi}\right)^{6/5} \times \left(\frac{M}{M_P}\right)^{8/5}, \tag{56}$$

which is a little larger than the total evaporated charge $\Delta Q \sim 10^{20}$. For pure leptonic flat direction such as LLe, Q balls must completely evaporate above the temperature where electroweak phase transition occurs in order to convert lepton number into baryon number by sphaleron effects, and hence $Q \gtrsim 10^{20}$ is not acceptable. On the other hand for flat directions carrying baryon number such as udd, $10^{20} \lesssim Q \lesssim 10^{22}$ is allowed. In such a case, where Q balls decay below the freeze-out temperature of LSP, we must care about the overproduction of LSPs from the Q-ball decay. But in our scenario entropy production from moduli decay dilutes them. Therefore for udd direction[4] the gauged $U(1)_{B-L}$ scenario in the presence of heavy moduli is marginally possible.

In Fig. 3, we show the resultant baryon asymmetry in the (M, v) plane with constraints. We can see that for $m_\chi = 100$ TeV, Q balls become too large. But for larger m_χ the correct baryon asymmetry can be obtained without forming too large Q balls.

[4] Actually LLe and udd direction can have large field value simultaneously.

V. AFFLECK-DINE BARYOGENESIS WITHOUT SUPERPOTENTIAL

A. The model

Next we consider the models of Affleck-Dine baryogenesis with gauged $U(1)_{B-L}$ including no nonrenormalizable superpotentials due to some symmetry such as R symmetry. In such a case, baryon number violating operators are supplied by higher order effects from Kahler potentials (see e.g., Ref. [41]) and the initial amplitude of the Affleck-Dine field can become as large as the Planck scale. The dynamics of the Affleck-Dine field is similar to the previous section. As a result, large Q balls are formed associated with Affleck-Dine baryogenesis and they decay at late time after the freezeout of LSPs but before BBN. Interestingly, in this type of model late-decaying Q balls may once dominate the Universe [41]. If this is the case, a nice feature arises when considering the moduli-induced gravitino problem, as explained in Sec. VI.

Now let us investigate the above model. The zero-temperature scalar potential for the flat direction ψ is given by

$$V(\psi) = (m_\psi^2 - c_H H^2)|\psi|^2 + \frac{m_{3/2}^2}{nM^{n-2}}(a_m\psi^n + \mathrm{H.c.}) + \frac{H^2}{nM^{n-2}}(a_H\psi^n + \mathrm{H.c.}) + \ldots, \tag{57}$$

where the ellipsis denote the higher order terms, which stabilize the Affleck-Dine field at some value of order the Planck scale. Note that the potential (57) also has charge and/or color breaking global minimum near the field value at M. Similar to the previous section, in order to avoid falling into this minimum, the D term stopping at v must satisfy the following condition,

$$v \lesssim |\psi|_{\mathrm{hill}} \sim \frac{m_\psi}{m_{3/2}}M. \tag{58}$$

Here we consider only the case without early oscillation due to thermal effects. This requires

$$T_R \lesssim \alpha^{-1}m_\psi^{1/2}|\psi_0|M_P^{-1/2} \sim 2\times 10^{11}\ \mathrm{GeV}\left(\frac{0.1}{\alpha}\right)\left(\frac{m_\psi}{100\ \mathrm{GeV}}\right)^{1/2}\left(\frac{|\psi_0|}{M_P}\right). \tag{59}$$

At the beginning of the oscillation $H = m_\psi$, the baryon number is calculated as

$$n_B(t_{\mathrm{os}}) \simeq \frac{|a_m|\delta_e m_{3/2}^2}{m_\psi M^{n-2}}\psi_0^n, \tag{60}$$

where ψ_0 is given by the $U(1)_{B-L}$ breaking scale v. Note that we assume the Affleck-Dine field has baryonic charge.

In such a case, the whole baryon number created by the coherent motion of the Affleck-Dine field contributes to the baryon number of the Universe as far as Q balls decay before BBN. If it does not have baryonic charge and only has leptonic charge, only some fraction of the total created lepton number evaporated from Q balls at the temperature above the electroweak scale can be converted into baryon number through the sphaleron effects [42,43]. Thus the resultant baryon asymmetry is suppressed. We do not consider such a case.

The charge of Q balls is given by Eq. (55) where the ellipticity parameter ϵ is now estimated as

$$\epsilon \sim \delta_e \left(\frac{m_{3/2}}{m_\psi}\right)^2 \left(\frac{v}{M}\right)^{n-2}.$$

Thus we obtain the charge of Q balls

$$Q \sim 4 \times 10^{26} \left(\frac{\gamma}{6\times 10^{-3}}\right)\left(\frac{1\ \text{TeV}}{m_\psi}\right)^2 \left(\frac{v}{M_P}\right)^2$$

for $\epsilon \lesssim 0.01$, and

$$Q \sim 4 \times 10^{32} \left(\frac{\gamma}{6\times 10^{-3}}\right)\left(\frac{m_{3/2}}{100\ \text{TeV}}\right)^2 \left(\frac{1\ \text{TeV}}{m_\psi}\right)^4 \left(\frac{v}{M_P}\right)^n \times \left(\frac{M_P}{M}\right)^{n-2}$$

for $\epsilon \gtrsim 0.01$. On the other hand, the decay temperature of Q balls is estimated as [22,44]

$$T_Q \sim 1.8\ \text{GeV} \left(\frac{m_\psi}{1\ \text{TeV}}\right)^{1/2} \left(\frac{10^{24}}{Q}\right)^{1/2}$$

if there exist lighter scalar fields than the Affleck-Dine field, and hence we can see that Q balls decay below the electroweak scale but before BBN for some parameter region, $v \ll M_P$ and/or $M \gg M_P$. The entire cosmological scenario depends on the decay temperature of Q balls T_Q, the initial amplitude of the Affleck-Dine field v, and that of the modulus field χ_0. We assume the reheating temperature from the inflaton is not so high as the inflaton dominates the Universe when the Affleck-Dine field begins to oscillate but decays well before the modulus field decays. This is satisfied for $10\ \text{GeV} \lesssim T_R \lesssim 10^9$ GeV. The following analysis does not depend on the precise value of T_R as far as the T_R lies in the above range.

The final reheating comes from moduli or Q balls. If the following condition

$$T_Q < T_\chi \left(\frac{v}{\chi_0}\right)^2 \tag{61}$$

is satisfied (T_χ is the decay temperature of the modulus field), Q balls dominate the Universe before they decay but after the modulus field decays.

B. Baryon asymmetry

In the case with Q-ball domination, the baryon-to-entropy ratio is fixed at the decay of Q balls,

$$\frac{n_B}{s} = \frac{n_B}{\rho_\psi}\frac{\rho_\psi}{s} = \frac{n_B}{\rho_\psi}\frac{3T_Q}{4} \tag{62}$$

and the subsequent cosmological scenario does not depend on the properties of moduli. This is calculated as

$$\frac{n_B}{s} \sim 7 \times 10^{-3} \left(\frac{6\times 10^{-3}}{\gamma}\right)^{1/2}\left(\frac{1\ \text{TeV}}{m_\psi}\right)^{3/2}\left(\frac{m_{3/2}}{100\ \text{TeV}}\right)^2 \times \left(\frac{M_P}{M}\right)^2\left(\frac{v}{M_P}\right)$$

for $\epsilon \lesssim 0.01$, and

$$\frac{n_B}{s} \sim 2 \times 10^{-5} \left(\frac{6\times 10^{-3}}{\gamma}\right)^{1/2}\left(\frac{1\ \text{TeV}}{m_\psi}\right)^{1/2}\left(\frac{m_{3/2}}{100\ \text{TeV}}\right) \times \left(\frac{M_P}{M}\right)^{(n-2)/2}\left(\frac{v}{M_P}\right)^{(n-4)/2},$$

for $\epsilon \gtrsim 0.01$, where we have assumed $\delta_e \sim 0.01$. On the other hand, in the case of no Q-ball domination, the final reheating occurs due to the modulus decay. The baryon-to-entropy ratio is thus given by

$$\frac{n_B}{s} = \frac{n_B}{\rho_\chi}\frac{\rho_\chi}{s} = \frac{n_B}{\rho_\chi}\frac{3T_\chi}{4} \tag{63}$$

and it is estimated as

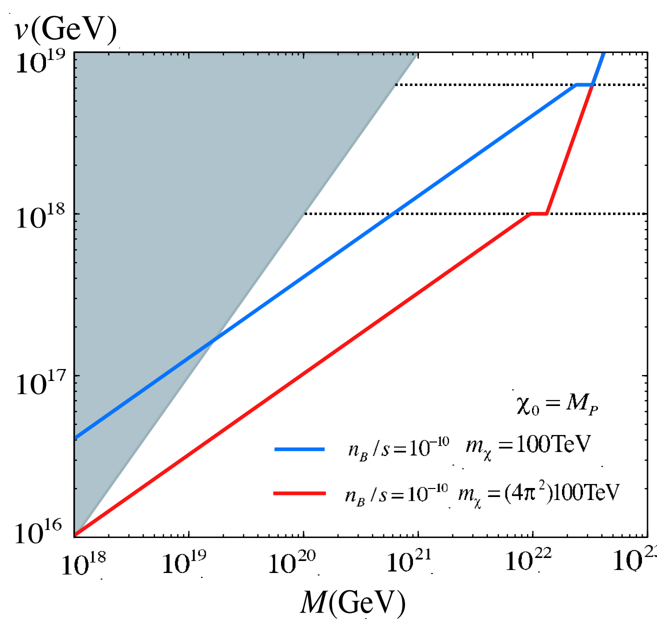

FIG. 4 (color online). The two solid lines show $n_B/s \sim 10^{-10}$ for $m_\chi = 100$ TeV and $(4\pi^2)100$ TeV in the no superpotential model. We take $m_{3/2} = 100$ TeV. In the dark shaded region the Affleck-Dine field is trapped into charge-breaking minima and baryogenesis does not work.

$$\frac{n_B}{s} \sim 2 \times 10^{-4}\sqrt{c}\left(\frac{m_\chi}{100\ \text{TeV}}\right)^{3/2}\left(\frac{1\ \text{TeV}}{m_\psi}\right)^3\left(\frac{m_{3/2}}{100\ \text{TeV}}\right)^2 \times \left(\frac{M_P}{M}\right)^{n-2}\left(\frac{v}{M_P}\right)^n\left(\frac{M_P}{\chi_0}\right)^2. \tag{64}$$

In both cases, we can see that $M \gg M_P$ and/or $v \ll M_P$ is required in order to obtain the correct order of baryon asymmetry.

We show in Fig. 4 the contour where the appropriate baryon asymmetry $n_B/s \sim 10^{-10}$ is obtained for $n = 4$ and the modulus mass $m_\chi = 100$ TeV and $(4\pi^2)100$ TeV. It should be noticed that for $v > \sqrt{3}M_P$, a brief period of inflation occurs due to the Affleck-Dine field. But in general, supergravity effects steepen the potential above the Planck scale, and hence the region with $v \gtrsim M_P$ is not favored from naturalness. Above the dotted lines Q-ball domination is realized for each modulus mass. It can be seen that for $M \gtrsim 10^{18}$ GeV and $v \gtrsim 10^{16}$ GeV, this baryogenesis mechanism works. But another subtlety arises when one considers the LSP produced by the Q-ball decay or gravitinos from modulus decay. This will be discussed in Sec. VI.

VI. REMARKS ON HEAVY MODULUS DECAY

Before closing the discussion, we consider some nontrivial features of the modulus decay. In the above arguments, we have not considered the details of the decay products of the modulus field. We briefly discuss the other consequences of modulus decay on cosmological evolution.

A. Nonthermal dark matter from modulus decay

One of the favored natures of the supersymmetric theory is that it can provide the candidate of the dark matter of the Universe. Under the R-parity conservation, the LSP becomes stable and if it has the appropriate annihilation cross section, it can account for the energy density of the dark matter [45]. In anomaly-mediated SUSY breaking models, winolike neutralino naturally becomes the LSP. In the standard thermal relic scenario, the mass of wino should be as heavy as 2 TeV to account for the dark matter because of its large annihilation cross section.[5] In the mirage-mediation model, the LSP is the mixed state of the bino and Higgsino-like neutralino [47,48] and their thermal relic abundance can account for the dark matter of the Universe [49]. But in our scenario, the final reheating temperature from modulus decay is typically much lower than the freeze-out temperature of the LSP and the thermal relic cannot be the dark matter. However, there is a nonthermal origin of the dark matter from the decay of moduli, and there arises a possibility that nonthermal LSPs can account for the dark matter of the Universe. Its abundance is estimated as [16,50]

$$Y(T) \simeq \left[\frac{1}{Y(T_\chi)} + \sqrt{\frac{8\pi^2 g_*}{45}}\langle\sigma v\rangle M_P(T_\chi - T)\right]^{-1}, \tag{65}$$

where $Y = n_{\rm LSP}/s$ and T_χ denotes the decay temperature of the modulus field. We can see that if the annihilation cross section is large enough, the second term dominates and the relic abundance is inversely proportional to it. In terms of the density parameter, we can rewrite it as

$$\Omega_{\rm LSP}h^2 \sim 0.27\left(\frac{10}{g_*(T_\chi)}\right)^{1/2}\left(\frac{100\ \text{MeV}}{T_\chi}\right)\left(\frac{m_{\rm LSP}}{100\ \text{GeV}}\right)^3 \times \left(\frac{10^{-3}}{m_{\rm LSP}^2\langle\sigma v\rangle}\right). \tag{66}$$

For the Higgsino-like neutralino LSP, the annihilation cross section into the W-boson pair is estimated as [51]

$$\langle\sigma v\rangle \simeq \frac{\pi\alpha_2^2}{2}\frac{m_{\rm LSP}^2}{(2m_{\rm LSP}^2 - m_W^2)^2}\left(1 - \frac{m_W^2}{m_{\rm LSP}^2}\right)^{3/2} \tag{67}$$

and for the winolike neutralino LSP, it is given by [16]

$$\langle\sigma v\rangle \simeq 8\pi\alpha_2^2\frac{m_{\rm LSP}^2}{(2m_{\rm LSP}^2 - m_W^2)^2}\left(1 - \frac{m_W^2}{m_{\rm LSP}^2}\right)^{3/2}, \tag{68}$$

where α_2 is the $SU(2)_L$ gauge coupling constant and m_W is the mass of the W boson. We can see that the desired density of the LSP can be obtained for $m_{\rm LSP} \sim 100$ GeV. Thus both the dark matter and baryon asymmetry of the Universe can be explained even in the presence of heavy modulus fields.

One may consider that these nonthermal LSPs from late-decaying particles can have large free-streaming length $\lambda_{\rm FS}$ ($\gtrsim 1$ Mpc) and may become the warm dark matter [15,52–54]. Now we estimate the free-streaming length of nonthermally produced LSPs assuming that they contribute to the dark matter of the Universe. For simplicity, we neglect the energy loss via the interaction between particles in the thermal bath. The free-streaming length of the LSP produced at τ_χ is given by [55]

$$\lambda_{\rm FS} \sim 1.0\ \text{Mpc}\, u_d\left(\frac{\tau_\chi}{10^6\ \text{sec}}\right)^{1/2} \times \left\{1 + 0.14\ln\left[\left(\frac{10^6\ \text{sec}}{\tau_\chi}\right)^{1/2}\frac{1}{u_d}\right]\right\}, \tag{69}$$

where $u_d = \sqrt{m_\chi^2 - 4m_{\rm LSP}^2}/2m_{\rm LSP}$. Using $\tau_\chi \sim (4\pi M_P^2/cm_\chi^3)$, we can rewrite it as

[5] There is an argument that the nonperturbative effect enhances the annihilation cross section, and wino should be about 3 TeV if its thermal relic accounts for the dark matter of the Universe [46].

$$\lambda_{\rm FS} \sim 1.7 \times 10^{-2}\ {\rm Mpc}c^{-1/2}\left(\frac{100\ {\rm TeV}}{m_\chi}\right)^{1/2}\left(\frac{1\ {\rm TeV}}{m_{\rm LSP}}\right) \times \left\{1 + 0.07 \ln\left[c^{1/2}\left(\frac{m_\chi}{100\ {\rm TeV}}\right)^{1/2}\left(\frac{m_{\rm LSP}}{1\ {\rm TeV}}\right)\right]\right\}. \tag{70}$$

Thus, for $m_\chi \gtrsim 100$ TeV, the free-streaming length is much smaller than 1 Mpc and nonthermal LSPs serve as the cold dark matter. Actually, there exist non-negligible interactions of LSPs with background particles. It is expected that LSPs lose their energy and momentum through those interactions and hence the nonthermal LSPs from modulus decay unlikely take a role of warm dark matter [56,57].

B. Gravitinos from modulus decay

If $m_\chi > 2m_{3/2}$, which is naturally realized in the mirage-mediation model, the modulus decay into two gravitinos is kinematically allowed. In particular, it is found that such a decay mode generally has the branching ratio as large as $O(0.01)$ [58,59] and the late-decay of gravitinos generated in this way may cause another cosmological difficulty. In our scenario these gravitinos do not upset BBN, since they are also heavy enough to decay before the beginning of BBN. But LSPs emitted from the decay of those nonthermally produced gravitinos may overclose the Universe.

The decay of such nonthermal gravitinos does not release a huge entropy, because the energy density of the gravitino is 2 orders of magnitude smaller than that of the ordinary radiation. We denote the temperature at the modulus decay and at the gravitino decay as T_χ and $T_{3/2}$, respectively. The branching ratios of moduli that decay into ordinary radiation and two gravitinos are denoted as $B_r \sim O(1)$ and $B_{3/2}$. In these terms, the ratio of the energy density of the gravitino to radiation at the decay of the gravitino is given by

$$\frac{\rho_{3/2}}{\rho_r} = \epsilon \frac{B_{3/2}}{B_r}, \tag{71}$$

where

$$\epsilon = \begin{cases} \frac{T_{\rm NR}}{T_{3/2}} & (T_{\rm NR} > T_{3/2}) \\ 1 & (T_{\rm NR} < T_{3/2}) \end{cases} \tag{72}$$

and $T_{\rm NR}(= (m_{3/2}/m_\chi)T_\chi)$ denotes the temperature at which gravitinos become nonrelativistic. Since $B_{3/2} \ll B_r$, $\rho_{3/2}$ is not larger than ρ_r for $T_{\rm NR} < T_{3/2}$. When the gravitino becomes nonrelativistic before decay, we obtain

$$\frac{\rho_{3/2}}{\rho_r} = \frac{B_{3/2}}{B_r}\frac{T_\chi}{T_{3/2}}\frac{m_{3/2}}{m_\chi} \sim \frac{B_{3/2}}{B_r}\left(\frac{m_\chi}{m_{3/2}}\right)^{1/2}. \tag{73}$$

This ratio does not exceed 1 in the parameter region we are interested in, and baryon asymmetry is not diluted further by the gravitino decay. The difficulty comes from the subsequent decay of gravitinos into LSPs. The LSP abundance emitted from the gravitino decay is also expressed by Eq. (65) after replacing T_χ with $T_{3/2}$. But for $m_{3/2} \sim 100$ TeV, $T_{3/2}$ is so small that the LSPs do not annihilate and their density overcloses the Universe. To avoid this difficulty, we require $m_{3/2} \gtrsim 10^3$ TeV. But in such a case, the mass of the LSP becomes too large in anomaly-mediation or mirage-mediation models and the overclosure problem of LSPs is not cured.

Here we describe some ways to avoid the LSP overproduction problem in such a heavy moduli scenario with $m_\chi \gg m_{3/2}$. One possible solution is to reduce $B_{3/2}$ so that the abundance of gravitinos from moduli decay can be neglected. Depending on the Kahler potential and SUSY breaking sector, the branching ratio of modulus decay into gravitinos may have the helicity suppression factor $\sim(m_{3/2}/m_\chi)^2$ compared with other decay modes [59].

Another is to introduce lighter R-odd particles other than MSSM particles. Axino, which appears in the supersymmetric extension of the axion models [60], is one of the candidates [61]. In such a case, the overproduced lightest neutralino eventually decays into axinos and its abundance is reduced by the factor $(m_{\tilde{a}}/m_{\rm LSP})$ where $m_{\tilde{a}}$ is the axino mass. If $m_{\tilde{a}}$ is sufficiently small [62], the overproduction problem of the neutralino LSP can be solved. It should be noticed that it may also open the decay mode of the gravitino into axino and axion [63], and these newly produced nonthermal axions serve as the additional radiation energy density [64], which speeds up the Hubble expansion and changes the result of BBN especially the ^{4}He abundance. In terms of the effective number of neutrinos N_ν, the success of BBN requires $\Delta N_\nu \lesssim 1$ at the beginning of BBN.[6] But in our situation this is not a problem, since the gravitino abundance is smaller than the radiation at the decay of gravitinos and hence the axion abundance generated from the gravitino decay is also smaller than the radiation energy density. In this scenario axinos should decay before BBN [66]. Otherwise, decay products of the neutralino spoils BBN. This requires the Peccei-Quinn scale 10^{10} GeV $\lesssim F_{\rm PQ} \lesssim 10^{11}$ GeV, where the lower bound comes from the astrophysical consideration [67] and the upper bound comes from the requirement that the lifetime of the decaying neutralino into axino should be shorter than 1 sec. Thermally produced axinos [68] can contribute to the only small fraction of the energy density of the Universe because the final reheating temperature is very low. The coherent oscillation of the axion is also diluted by the modulus decay and has neglecting effects on cosmology for $F_{\rm PQ} \lesssim 10^{11}$ GeV [69].

Finally, we mention the possibility that gravitinos are diluted by entropy production after the modulus decay in

[6]The recent analysis of the primordial ^{4}He abundance favors the nonstandard value of $N_\nu(>3)$ [65], but $\Delta N_\nu \sim 1$ is disfavored.

the following subsection. This is already built in the models of the Q-ball dominant universe (Sec. V), as we will see.

C. Dilution by Q-ball decay

The abundance of gravitinos from modulus decay is expressed as

$$Y_{3/2} = 2\frac{B_{3/2}}{B_r}\frac{3T_\chi}{4m_\chi}. \tag{74}$$

Assuming no annihilation, the energy density of the LSP produced by the decay of gravitinos is given by

$$\Omega_{\rm LSP}h^2 \sim 2.3 \times 10\sqrt{c}\left(\frac{B_{3/2}/B_r}{0.01}\right)\left(\frac{m_{\rm LSP}}{100\ {\rm GeV}}\right)\left(\frac{m_\chi}{100\ {\rm TeV}}\right)^{1/2}. \tag{75}$$

Thus we need the dilution factor $\Delta \sim 10^2$–10^3 after the production of gravitinos. In fact, in the Q-ball dominant case investigated in Sec. V, the decay of Q balls releases large entropy and dilutes the gravitinos from modulus decay. If we denote the decay temperature of Q balls as T_Q, the dilution factor is given by

$$\Delta = \frac{T_\chi}{T_Q}\left(\frac{|\psi_0|}{|\chi_0|}\right)^2. \tag{76}$$

Thus if T_Q is slightly smaller than T_χ, and the initial amplitude of the modulus field $|\chi_0|$ is slightly smaller than that of the Affleck-Dine field $|\psi_0|$, the required dilution factor is obtained. This is no more than the situation we encountered in Sec. V, and hence those Q balls give the desired dilution factor, solving the overproduction problem of LSPs from gravitino decay. Note also that LSPs are also produced from the decay of Q balls. Their abundance is given by Eq. (65) after replacing T_χ with T_Q. As far as $T_Q \gtrsim 100$ MeV, LSPs can effectively annihilate and their abundance becomes below or comparable to that of the dark matter [17,50].

In Fig. 5, the result with $|\chi_0| = 10^{17}$ GeV is shown. It can be seen that the wide parameter region which has been favored in Sec. V is excluded by the constraint from the overproduction of LSPs both from the Q-ball decay (the purple shaded region) and gravitino decay (the blue shaded region). We can see that only in the narrow parameter region for $M \sim 10^{22}$ GeV and $v \sim M_P$, both the dark matter as nonthermal LSPs and baryon asymmetry can be explained simultaneously. As the value of χ_0 is reduced, the constraint is relaxed. Although some degree of tunings to the parameters, especially the initial amplitude of the modulus and Affleck-Dine field, is required, this is a fully consistent cosmological scenario in the presence of heavy moduli.

Finally we comment on the possibility where the similar dilution from the Q-ball decay is obtained in the model of Sec. III. In order to realize this, we consider the situation where a flat direction other than the Affleck-Dine field

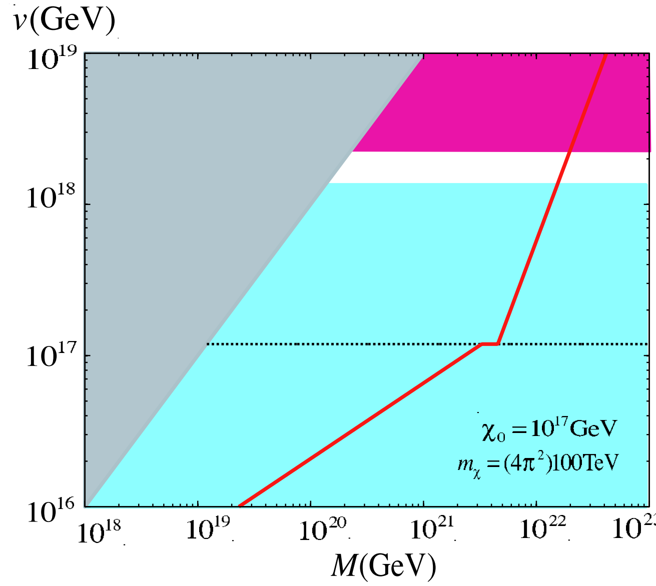

FIG. 5 (color online). Same as Fig. 4, but for $\chi_0 = 10^{17}$ GeV. We also show the constraint from the LSP overproduction from gravitino decay for $m_\chi = 4\pi^2 100$ TeV. In the lower right blue shaded region, the Q-ball decay cannot dilute LSPs produced by the decay of moduli-induced gravitinos sufficiently. In the upper right purple shaded region, LSPs produced by the Q-ball decay overclose the Universe.

responsible for baryogenesis dominates the Universe after the moduli decay. In fact, udd and LLe directions can have the large field value simultaneously. As described in Sec. III, the $n = 6$ udd direction is used as the Affleck-Dine field which creates the appropriate baryon number. On the other hand, the LLe direction also has the large field value. We assume for the LLe direction there does not exist nonrenormalizable superpotentials which lift the direction and we parametrize this direction as ψ (such a model was considered in Ref. [70]). Similar to the case in Sec. V, ψ has the initial amplitude ψ_0 of order M_P. The late decay of Q balls from the ψ-condensate dilutes the gravitino to the cosmologically safe value. Although it also dilutes the baryon asymmetry by the factor $\Delta \sim 10^2$, dilution of such an amount of baryon asymmetry is not so harmful, as can be seen from Fig. 2.

VII. CONCLUSIONS

Within the framework of the fundamental theory such as supergravity or superstring theory, there appear cosmologically harmful scalar fields called moduli. In anomaly-mediated SUSY breaking or the mirage-mediation model, moduli are heavy and decay well before BBN starts, but the decay process dilutes the preexisting baryon asymmetry. We have shown that in some models of the Affleck-Dine baryogenesis mechanism, a large amount of baryon asymmetry can be generated and can survive the dilution from the modulus decay. Successful baryogenesis requires high reheating temperature from the inflaton, $T_R \gtrsim 10^{10}$ GeV, and high effective cutoff scale, $M \gtrsim 10^{19}$ GeV for early oscillation models. Such a high reheating temperature is naturally realized in chaotic inflation models [71]. Gauged $\mathrm{U}(1)_{B-L}$ models also work for some parameter regions. We also investigated the gauged $\mathrm{U}(1)_{B-L}$ model without superpotentials which lift the flat direction. The favored parameter region is also found in this type of model. Other baryogenesis mechanisms such as thermal leptogenesis [72] and electroweak baryogenesis [73] do not work, since produced baryon asymmetry is not so large as to survive the dilution.

Aside from baryon asymmetry, dark matter of the Universe can also be explained by the nonthermal LSPs from the decay of moduli. The final reheating temperature is determined by the decay of moduli and it is predicted as from a few MeV to 1 GeV for 100 TeV $\lesssim m_\chi \lesssim (4\pi^2)100$ TeV. Hence the standard cosmological scenario below a few MeV should not be changed. One subtlety arises when we consider the gravitino production from decay of the heavy moduli if the modulus mass m_χ is larger than 2 times the gravitino mass $m_{3/2}$. If the branching ratio of modulus decay into two gravitinos is not suppressed, we encounter another cosmological problem, i.e., the overproduction of neutralino LSPs from the subsequent decay of gravitinos. In the Q-ball dominant scenario in Sec. V, Q-ball decay dilutes the gravitino and the problem can be solved by choosing the initial amplitude of the modulus and the Affleck-Dine field as $|\chi_0| \lesssim |\psi_0| \sim M_P$. Besides the Q-ball dominant scenario, there are a couple of ways to avoid this difficulty. One is controlling the SUSY breaking sector to suppress the branching ratio into gravitinos, and another is to introduce the axinos. The other solution is to invoke another flat direction condensate into large Q balls. The late decay of Q balls dilutes the gravitino abundance, and also such a Q-ball decay itself can provide the nonthermal origin of the dark matter, similar to the Q-ball dominant model. In any way, our scenario provides the realistic cosmological scenario in the presence of modulus fields and may have a phenomenologically interesting implication to future collider experiments and direct or indirect detection of the dark matter.

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