Affleck-Dine baryogenesis and leptogenesis with a gauged $U(1)_{B-L}$

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We briefly review the present status of Affleck-Dine baryogenesis and leptogenesis scenarios in the minimal supersymmetric standard model (MSSM) in the context of the gravity-mediated supersymmetry (SUSY) breaking, and show that there is a serious cosmological problem in the Affleck-Dine mechanism. That is, the late decay of the associated large Q balls leads to the over production of the lightest supersymmetric particles. Then we point out that the minimal extension of the MSSM by introducing a gauged $U(1)_{B-L}$ symmetry naturally solves this problem. Here, the breaking scale of the $U(1)_{B-L}$ can be determined quite independently of the reheating temperature from the required baryon asymmetry. It is extremely interesting that the obtained scale of the $U(1)_{B-L}$ breaking is quite consistent with the one suggested from the seesaw mechanism to explain the recent neutrino-oscillation experiments. We consider that the present scenario provides a new determination of the $U(1)_{B-L}$ breaking scale fully independent of the neutrinos masses. We also comment on the viability of the present scenario in anomaly mediated SUSY breaking models.

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I. INTRODUCTION

The Affleck-Dine (AD) mechanism [1,2] is one of the interesting possibilities for generating the present baryon (matter-antimatter) asymmetry in the early universe. This is expected to work naturally in the supersymmetric (SUSY) standard model,[1] since it has a number of flat directions (complex scalar fields) [5] carrying baryon (B) or lepton (L) charges. A linear combination of squark and/or slepton fields, which we will call an AD field, may have a large expectation value along a flat direction during an inflationary stage in the early universe. The AD field starts its coherent oscillation after the inflation ends and it creates a large net baryon and/or lepton asymmetry, which is finally transferred to matter particles when it eventually decays.

In recent developments, however, there appear several serious obstacles to the Affleck-Dine mechanism. In particular, the formation of a Q ball, which is a kind of a nontopological soliton [6,7], is the most serious. The Q balls are likely formed due to spatial instabilities of the coherent oscillation of the AD field [8]. It has been, in fact, shown in detailed numerical calculations that almost all the initial charges which the AD field carries are absorbed into the formed Q balls [9,10]. Thus, the resultant baryon asymmetry must be provided by decay of the associated Q balls, not of the AD field. These Q balls have very significant consequences on Affleck-Dine baryogenesis and cosmology [11,12].

To show how the difficulty arises from the Q-ball formation in the Affleck-Dine mechanism, let us briefly summarize the present status of the Affleck-Dine baryogenesis/leptogenesis scenarios in the context of the minimal SUSY standard model (MSSM) with R-parity conservation. The essential ingredients for the Affleck-Dine mechanism to determine the resultant baryon asymmetry are the initial amplitude of the AD field (more precisely, the amplitude when it starts the coherent oscillation) and the size of operators which kick this condensate to give it a phase rotational motion. An important point in this regard is that during the inflation and in the epoch dominated by oscillations of the inflaton, a large energy density of the universe violates SUSY, which induces a SUSY-breaking mass term of the order of the Hubble parameter H [2] (we will call it a Hubble-mass term) for the AD field. In the cases where the AD field has the minimal Kähler potential, the induced Hubble-mass term is positive, which drives the AD field towards the origin during the inflation, and hence the Affleck-Dine mechanism does not work. However, in general cases where the AD field has nonminimal couplings to the inflaton field, the induced Hubble-mass term can be negative so that the AD field develops a large expectation value. This is crucial for the Affleck-Dine mechanism to work and we assume this is the case in the following discussion.

The initial amplitude and evolution of the AD field are determined in a balance between the induced negative Hubble-mass term and nonrenormalizable operators which lift the associated flat direction. In the case where all of the nonrenormalizable operators in the superpotential are forbidden by some chiral symmetries, such as R symmetry, the initial amplitude of the AD field is naturally expected to be the reduced Planck scale $M_* \simeq 2.4 \times 10^{18}$ GeV which is the balance point between the induced negative Hubble-mass term and nonrenormalizable interactions in the Kähler potential. The operators which kick the AD field to give it a phase rotational motion are also provided by nonrenormalizable interactions in the Kähler potential. This is the simplest and a

[1] In most part of this paper, we assume the gravity-mediated SUSY breaking. We note on the Affleck-Dine mechanism in the anomaly-mediated SUSY breaking [3,4] in the last section.

quite possible scenario, but it generally predicts too much an amount of baryon asymmetry. Furthermore, even if it is diluted by late-time entropy productions, the decay of the associated Q balls causes a serious problem in the universe. This is because the decay temperature of the Q ball, which is inversely proportional to the square root of its charge, is expected to be well below the freeze-out temperature of the lightest SUSY particle (LSP),[2] which invalidates this scenario by means of the following argument.

If the decay temperature of the associated Q ball is lower than the freeze-out temperature of the LSP, the resultant baryon asymmetry and the abundance of the LSP cold dark matter is related by the following equation;

$$\Omega_\chi = 3\left(\frac{N_\chi}{3}\right) f_B \left(\frac{m_\chi}{m_n}\right) \Omega_B , \tag{1}$$

where N_χ is the number of LSP's produced per one baryon number, which is at least 3, $f_B \simeq 1$ is the fraction of baryon number stored in the form of Q balls, and m_n and m_χ are the nucleon mass and the neutralino (χ) LSP mass, respectively. Ω_X denotes the ratio of the energy density of X to the critical density of the present universe. Here we assume the case of the standard B-ino–like LSP.[3] Using the conservative constraint on the baryon asymmetry from the big-bang nucleosynthesis, $0.004 \lesssim \Omega_B h^2 \lesssim 0.023$ [14], the above relation leads to a stringent constraint on the neutralino LSP mass as

$$m_\chi \lesssim 16 \ \text{GeV} \left(\frac{h}{0.7}\right)^2 \left(\frac{\Omega_\chi}{0.4}\right) \left(\frac{3}{N_\chi}\right) \left(\frac{1}{f_B}\right). \tag{2}$$

The direct experimental lower bound on the neutralino mass is $m_\chi > 32.3$ GeV [15], and it excludes the above scenario. This is an inevitable consequence of a variety of Affleck-Dine baryogenesis scenarios[4] with the formation of large Q balls. Note that the Q balls are formed when $H \sim \mathcal{O}(\text{GeV})$,[5] which is often the epoch dominated by the oscillating inflaton, and the entropy productions which occur after this epoch cannot change the size of the Q balls, and hence their decay temperature.

On the other hand, if the general nonrenormalizable superpotential which is consistent with the MSSM gauge symmetries and R parity is present, most of the flat directions are lifted by dimension six operators which are F-term contributions from the dimension four superpotential [5], and their phase rotational motions are caused by the associated A terms. Unfortunately, however, all of these A terms, with only one interesting exception using the LH_u flat direction [16], preserve $B-L$ and the resultant B and L asymmetry are washed out by the ''sphaleron'' effects [17]. Thus, the baryon asymmetry in the present universe cannot be explained.[6] The almost unique viable scenario among the cases in which the initial amplitude of the AD field is fixed by the potential wall of the dimension six operators is the Affleck-Dine leptogenesis [16] using the LH_u flat direction.[7] This scenario has some particularly interesting features. The resultant baryon asymmetry is almost independent of the reheating temperature of inflation [19,20] and it is determined by only one parameter, i.e., the lightest neutrino mass, which results in an important prediction on the rate of the neutrinoless double beta decay [20].

Only one remaining possibility in the presence of general nonrenormalizable superpotential is to use the flat directions which can escape from the small window of the potential wall of the dimension six operators. In this case they are lifted by dimension ten operators which are F-term contributions of the dimension six superpotential. Here, all of the associated A terms violate $B-L$ and thus the present baryon asymmetry might be accounted for by using these flat directions. Unfortunately, this is not a viable scenario in the MSSM. The effective cutoff scale governing the dimension six operators in the superpotential is naturally considered to be the reduced Planck scale or higher. In this case the associated Q balls are too large to decay above the freeze-out temperature of the LSP [21]. Even if we extend the MSSM into the SUSY $SU(5)$ grand unified theories (GUT), this problem cannot be avoided. This is because all of the relevant flat directions are singlets under the GUT gauge group and the initial amplitude of the AD field (hence the size of the Q balls), cannot be suppressed.

In this paper, we consider a possibility to extend the MSSM gauge group. The most likely and minimal extension of the gauge group is to introduce the $U(1)_{B-L}$ symmetry which is supposed to be broken at a very high-energy scale [22]. If this is the case, anomaly cancellation conditions automatically require the existence of three right-handed neutrino chiral superfields which are singlets under the MSSM gauge group and carry $B-L$ charge $+1$. One can easily imagine that these right-handed neutrinos acquire large Majorana masses of the order of the $U(1)_{B-L}$ breaking scale. As a big bonus, we can naturally obtain a realistic mass spectrum for the lighter neutrinos via the so-called seesaw mechanism [23]. We stress, as we will see later, that this minimal and phenomenologically desirable extension of the MSSM indeed cures the Affleck-Dine baryogenesis/leptogenesis scenarios.

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[2]See the discussion in Sec. III and Appendix.

[3]As for the case of Higgino-like and W-ino–like LSP which have large annihilation cross sections, see Ref. [13].

[4]If we use purely leptonic flat directions such as $LL\bar{e}$, the associated leptonic Q balls (L balls) must evaporate above the electroweak scale to make the ''sphalerons'' effectively work, which results in a more stringent constraint. See the discussion in Sec. III.

[5]See the discussion in Sec. III.

[6]It might be possible to make large Q balls in order to protect the produced baryon asymmetry from the ''sphaleron'' effects. We find that it requires the existence of an unnaturally large cutoff scale $M \sim 10^{25}$ GeV. Furthermore, the size of the Q balls and the reheating temperature of inflation should be fine tuned so that one can avoid the above-mentioned LSP over-production problem.

[7]In this scenario, the associated L balls are expected to evaporate completely soon after their formation, or even not to be formed due to the strong thermal effects. Even if the the AD field starts its coherent oscillation when it is decoupled from the surrounding plasma, L balls are not formed or are very small because of the particular behaviors of the soft SUSY-breaking mass of this flat direction under renormalization-group equations [11,18].

In the following part of this paper, we perform an analysis of the Affleck-Dine baryogenesis/leptogenesis scenario which uses the dimension six superpotential in the context of the minimal extension of the MSSM with a gauged $U(1)_{B-L}$ symmetry. Surprisingly enough, all of the relevant flat directions which can pass through the small window of the potential wall coming from the dimension four superpotential have nonzero $B-L$ charges with the same sign. Thus, the D-term contribution of the $U(1)_{B-L}$ may stop the AD field to run toward the Planck scale before the dimension ten operators do so. This is a crucial point for our scenario. We find that the baryon asymmetry required from the big-bang nucleosynthesis can be naturally obtained with relatively low reheating temperatures enough to avoid a cosmological gravitino problem [24]. Furthermore, the initial amplitude of AD field suppressed by the $U(1)_{B-L}$ D-term contribution makes the expected Q balls small enough to decay well above the freeze-out temperature of the LSP. We also emphasize that the breaking scale of the $U(1)_{B-L}$ is almost uniquely determined to produce the required baryon asymmetry. We find the breaking scale $\simeq(2-7)\times10^{14}$ GeV. Note that the $U(1)_{B-L}$ breaking scale is determined totally independently of the neutrino masses. It is very encouraging that the obtained scale is quite consistent with the scale of the right-handed neutrino masses suggested from the seesaw mechanism to explain the recent neutrino-oscillation experiments [25].

II. THE SCENARIO FOR BARYOGENESIS AND LEPTOGENESIS

In this paper, we consider a minimal extension of the MSSM in which the gauge group is extended by introducing only $U(1)_{B-L}$, and assume this gauge symmetry is spontaneously broken at a very high-energy scale. We postulate that all renormalizable and nonrenormalizable operators allowed by gauge symmetries and R parity exist in the superpotential. The relevant flat directions which can pass through the small chinks of the potential wall of the dimension six operators are labeled by the following monomials of chiral superfields [5],

$$\overline{udd},\ LL\bar{e},\ \text{and a linear combination of } (\overline{udd}, LL\overline{ddd}). \tag{3}$$

Here, we mean that the flat directions labeled by $(\overline{udd},LL\overline{ddd})$ can pass through the dimension six potential wall with $\langle\overline{uddLL\overline{ddd}}\rangle\neq0$, but $\overline{udd}$ and $LL\bar{e}$ can do so only when $\langle\overline{uddLL\bar{e}}\rangle=0$. From Eq. (3), we can see that all of the flat directions carry nonzero $B-L$ charges with the same sign. Thus, the $U(1)_{B-L}$ D-term contribution from the relevant flat directions cannot be canceled out within themselves. Consequently, if the vacuum expectation values of some fields to break the $U(1)_{B-L}$ symmetry are stabilized during the inflation and in the inflaton-dominated epoch, these flat directions can be lifted at the breaking scale of the $U(1)_{B-L}$ as shown later. This is a crucial point for our scenario. We also comment here that these flat directions are $SU(5)_{\rm GUT}$ singlets, and hence the $SU(5)_{\rm GUT}$ D-term contribution cannot lift them.

First, let us show the evolution of the AD fields during the inflation. Here, we see that their initial values are really of the order of the breaking scale of the $U(1)_{B-L}$. In the following discussion, all of the superfields and their scalar components which parametrize the relevant flat directions will be represented symbolically by ϕ and treated as just a single field. This is not accurate in the case where there exist multiple flat directions. However, it does not alter the following order estimation of the baryon asymmetry and hence we use this representation for simplicity.

If the Kähler potential for the associated flat directions is in the minimal form (i.e., $K=\phi^\dagger\phi$), the induced SUSY-breaking mass for the AD field is positive as emphasized in Ref. [2], and is given by $\delta V\simeq3H^2|\phi|^2$. As described in the Introduction, in order to have a large expectation value of the AD field, nonminimal couplings for the AD field to the inflaton are required;

$$K=\phi^\dagger\phi+I^\dagger I+\frac{b_\phi}{M_*^2}\phi^\dagger\phi I^\dagger I+\cdots, \tag{4}$$

where I denotes the inflaton superfield.[8] The induced SUSY-breaking term from the couplings in Eq. (4) is approximately given by

$$\delta V\simeq3(1-b_\phi)H^2|\phi|^2, \tag{5}$$

and we assume $3(1-b_\phi)\sim-1$ for simplicity. In contrast to the ordinary pictures described in the introduction, in our scenario, the initial amplitude of the AD field is not determined in the balance between the negative Hubble-mass term and nonrenormalizable operators. To demonstrate our point, let us consider, for example, the following superpotential;

$$W=\lambda X(S\bar{S}-v^2), \tag{6}$$

where $\lambda=\mathcal{O}(1)$ is a coupling constant, v is the breaking scale of the $U(1)_{B-L}$, and X,S, and $\bar{S}$ are chiral superfields which are singlets under the MSSM gauge group and carry $0,+2,-2$ of $B-L$ charges, respectively. Then, the relevant scalar potential to determine the initial amplitude of the AD field is given by

$$V_D\simeq\frac{1}{2}g^2\left(2|S|^2-2\frac{v^4}{|S|^2}-|q||\phi|^2\right)^2-H^2|\phi|^2$$
$$+3(1-b_S)H^2|S|^2+3(1-b_{\bar{S}})H^2\frac{v^4}{|S|^2}, \tag{7}$$

[8]There might be three-point couplings of the inflaton to the AD and other fields in the Kähler potential, however, they do not alter the following discussion and we ignore them for simplicity.

where $g=\mathcal{O}(1)$ is the gauge coupling constant of the $U(1)_{B-L}$, $q(<0)$ is the $B-L$ charge of the AD field, and $b_S, b_{\bar{S}}$ are the nonminimal couplings for the $S, \bar{S}$ fields to the inflaton corresponding to the coupling b_ϕ in Eq. (4). In Eq. (7), we have eliminated the $\bar{S}$ field by minimizing the F term from the superpotential in Eq. (6); $|\bar{S}| \simeq v^2/|S|$, which can be justified as long as the $U(1)_{B-L}$ breaking scale v is much larger than the Hubble parameter during the inflation, $H_{\rm inf}$. From Eq. (7), one can show that the minimum for $|S|$ is given by

$$|S|^2 \simeq \frac{|q|}{4}|\phi|^2 + \sqrt{v^4 + \left(\frac{|q|}{4}\right)^2 |\phi|^4} \tag{8}$$

and the S and $\bar{S}$ fields have masses $\simeq v \gg H_{\rm inf}$ around this minimum. Thus, the S and $\bar{S}$ fields faithfully track this minimum throughout the history of the universe, and the contribution to the potential for the AD field from the first term in Eq. (7) practically vanishes [16]. Then, from Eqs. (7) and (8), the relevant potential for the AD field in the initial stage is given by

$$V_D(\phi) \simeq \begin{cases} \left((1-b_S)\dfrac{3|q|}{2} - 1\right) H^2|\phi|^2 + \mathcal{O}\left[H^2 v^4 / \left(\dfrac{|q|}{4}|\phi|^2\right)\right] \\ \quad \text{for } \left(\dfrac{|q|}{4}|\phi|^2 > v^2\right), \\ -\left(1 + (b_S - b_{\bar{S}})\dfrac{3|q|}{4}\right) H^2|\phi|^2 + \mathcal{O}\left[\dfrac{H^2}{v^2}\left(\dfrac{|q|}{4}|\phi|^2\right)^2\right] \\ \quad \text{for } \left(\dfrac{|q|}{4}|\phi|^2 < v^2\right). \end{cases} \tag{9}$$

Thus, we see from Eq. (9) that the potential minimum of the AD field and also those of the S and $\bar{S}$ are, in fact, of the order of the v during the inflation, if the coupling constants b_S and $b_{\bar{S}}$ satisfy the following conditions:

$$b_S \lesssim -1 \quad \text{and} \quad b_{\bar{S}} - b_S \lesssim 4, \tag{10}$$

where we have used the fact that $q \simeq -1/3$. During the inflation, the AD field evolves exponentially to this minimum since the effective mass of the AD field is of the order of the Hubble parameter.

The evolution of the AD field after the inflation is very simple. The AD field is frozen at this initial point $|\phi|_0 \simeq v$ until the Hubble parameter becomes smaller than the soft mass for the AD field in the true vacuum, and then, it starts the coherent oscillation. This is because the initial point is determined almost independently of the Hubble parameter, as easily seen from Eq. (7). Note that a similar argument is always possible as long as the superfields to break the $U(1)_{B-L}$ have masses much lager than the Hubble parameter, and hence the above stabilization mechanism of the AD field is not restricted to the specific form of the superpotential as in Eq. (6).

Here, let us make a brief comment on the effects of the nonrenormalizable couplings of the AD, S, and $\bar{S}$ fields to the inflaton in the Kähler potential. As can be seen from the discussion just described before, energy density coming from the AD, S, and $\bar{S}$ fields is $\mathcal{O}(H^2 v^2)$ and much smaller than that of the inflaton; $\mathcal{O}(H^2 M_*^2)$. Therefore such couplings do not alter the dynamics of the inflaton, if there are no couplings between the inflaton and the AD, S, and $\bar{S}$ fields in the superpotential. Furthermore, such nonminimal couplings in the Kähler potential are also desirable for a successful reheating.

Now, let us estimate the baryon asymmetry. The nonrenormalizable operators in the superpotential which provide the relevant A terms are given by

$$W = \frac{\lambda_1}{M_*^3}\left(\frac{S}{M_*}\right)(\overline{udd})^2, \quad \frac{\lambda_2}{M_*^3}\left(\frac{S}{M_*}\right)(LL\bar{e})^2, \tag{11}$$

where λ_1 and λ_2 are $\mathcal{O}(1)$ coupling constants. One may wonder why these $B-L$ conserving operators can produce nonzero $B-L$ charges, however, they actually do. It is clear from the above arguments on the evolutions of the AD and S fields that the S field sits at the $U(1)_{B-L}$ breaking scale v throughout the history of the universe, and hence the $U(1)_{B-L}$ symmetry is in fact spontaneously broken below the scale of v, where the AD field behaves nontrivially to produce present baryon asymmetry. In the following discussion, we express these superpotential operators as

$$W = \frac{1}{6M^3}\left(\frac{S}{M_*}\right)\phi^6, \tag{12}$$

for simplicity, where $M(\gtrsim M_*)$ is the effective cutoff scale. Then, the scalar potential for the AD field is given by

$$V = m_\phi^2|\phi|^2 + \frac{m_{3/2}}{6M^3}\left(\frac{S}{M_*}\right)(a_m\phi^6 + \text{H.c.}) + \frac{1}{M^6}\left(\frac{|S|}{M_*}\right)^2|\phi|^{10} + \frac{|\phi|^{12}}{36M^6M_*^2} + V_D, \tag{13}$$

where $|a_m| = \mathcal{O}(1)$, and m_ϕ is the effective soft mass for the AD field which is expected to be of the order of the gravitino mass $m_{3/2}$. There might exist another A term induced by the energy density of the universe dominated by the oscillating inflaton [2],

$$\frac{H}{6M^3}\left(\frac{S}{M_*}\right)(a_H\phi^6 + \text{H.c}), \tag{14}$$

where $|a_H| = \mathcal{O}(1)$. We can safely neglect this term since it does not alter the order of the final baryon asymmetry.[9] We require the initial amplitude of the AD field, $|\phi|_0$, and that of the S field, $|S|_0$, satisfy the condition

$$|\phi|_0 \lesssim \left(m_\phi M^3 \frac{M_*}{|S|_0} \right)^{1/4}, \quad (m_\phi M^3 M_*)^{1/5}. \tag{15}$$

Then, the third and fourth terms in Eq. (13) are always smaller than the last one V_D. Thus, these dimension 10 and 12 operators do not play any role in our scenario. If the condition in Eq. (15) is not satisfied, the resultant baryon or lepton number density when the Q balls are formed is larger than that in the case without a gauged $U(1)_{B-L}$ symmetry, and hence the LSP over-production problem becomes even worse.[10]

The evolution of the AD field is described by the following equation.

$$\ddot{\phi} + 3H\dot{\phi} + \frac{\partial V}{\partial \phi^\dagger} = 0. \tag{16}$$

The baryon or lepton number density is related to the number density of the AD field as

$$n = \beta i (\dot{\phi}^\dagger \phi - \phi^\dagger \dot{\phi}), \tag{17}$$

where β is the baryon or lepton charge carried by the AD field. From Eqs. (16) and (17), the evolution of the baryon or lepton number density is given by the following equation:

$$\dot{n} + 3Hn = 2\beta \mathrm{Im}\left(\frac{m_{3/2}}{M^3} \frac{S}{M_*} (a_m \phi^6) \right). \tag{18}$$

By integrating Eq. (18), we can obtain the resultant baryon or lepton number density at time t as

$$[R^3 n](t) = 2\beta \frac{m_{3/2}}{M^3 M_*} \int^t dt\, R^3 \mathrm{Im}(a_m S \phi^6), \tag{19}$$

where R denotes the scale factor of the universe. From Eq. (19), one can easily see that the production of the baryon or lepton number effectively terminates as soon as the AD field ϕ starts its coherent oscillations around the origin, since the integrand of the right-hand side of Eq. (19) decreases as fast as $\propto t^{-4}$. Then, at this time which is denoted by the subscript 0, the baryon or lepton number density is given by

$$n(t_0) \simeq \frac{4\beta}{9} \frac{m_{3/2}}{H_0} \frac{|a_m||S\phi^6|_0}{M_* M^3} \delta_{eff}, \tag{20}$$

where $\delta_{eff} = \sin[\arg(a_m) + \arg(S) + 6\ \arg(\phi)]$ represents an effective CP violating phase and is expected to be $\mathcal{O}(1)$ unless the initial phase of the AD field is fine tuned. Then, the ratio of the baryon or lepton number density[11] to the entropy density after the reheating process of the inflation is given by

$$\begin{aligned} \frac{n}{s} &\simeq \frac{\beta}{9} |a_m| \delta_{eff} \frac{T_R m_{3/2}}{H_0^3} \frac{|S\phi^6|_0}{M_*^3 M^3} \sim 2 \times 10^{-10} \left(\frac{T_R}{10^5\ \mathrm{GeV}} \right) \\ &\times \left(\frac{1\ \mathrm{TeV}}{m_\phi} \right)^2 \left(\frac{v}{5 \times 10^{14}\ \mathrm{GeV}} \right)^7 \left(\frac{m_{3/2}}{m_\phi} \right) \\ &\times \left(\frac{M_*}{M} \right)^3 \left(\frac{|S\phi^6|_0}{v^7} \right), \end{aligned} \tag{21}$$

where T_R is the reheating temperature of the inflation, and at the second line, we have used $|a_m| \simeq \delta_{eff} \simeq 1$, $H_0 \simeq m_\phi$. This ratio stays constant unless additional entropy productions take place. Here, we stress that the breaking scale of the $U(1)_{B-L}$ can be determined as $v \simeq (2-7) \times 10^{14}$ GeV almost independently of the reheating temperature of the inflation for 10^3 GeV $\lesssim T_R \lesssim 10^7$ GeV (This parameter region for the reheating temperature T_R should be taken to avoid the LSP over production and the thermal effects, as we will see below.) We can see that the initial values of the AD and S fields satisfy the condition in Eq. (15) with $|\phi|_0, |S|_0 \sim v$ and $M \gtrsim M_*$. One should note that, in the second line of Eq. (21), we have assumed the absence of the relevant thermal effects which cause the early oscillation of the AD field. If the induced thermal effects dominate the effective potential for the AD field and cause the early oscillation, i.e., $H_0 > m_\phi$, the resultant baryon/lepton asymmetry is strongly suppressed [27], as easily seen from the first line of Eq. (21).

First, let us investigate the conditions to avoid the early oscillation due to the thermal mass terms. The field which couples to the AD field through the coupling constant f_i gets an effective mass $f_i|\phi|$. If this effective mass is smaller than the temperature T, the thermal fluctuations of this field produces the thermal mass term for ϕ, which is given by $c_i^2 f_i^2 T^2 |\phi|^2$, where c_i is a $\mathcal{O}(1)$ constant. If this thermal mass $c_i f_i T$ exceeds the Hubble parameter in the regime $H > m_\phi$, it causes the early oscillation of the AD field. Thus, to avoid the early oscillation, the following two conditions should not be satisfied simultaneously during the regime $H > m_\phi$;

[9] See the related discussion in Ref. [19].

[10] In the cases without a gauged $U(1)_{B-L}$ symmetry, the baryon or lepton number density is $\propto M^{3/2}$. If the condition in Eq. (15) is not satisfied, the resultant baryon or lepton number density when the Q balls are formed is enhanced by the factor $(M_*/|S|_0)^{1/2} \gg 1$, compared with that in the case without a gauged $U(1)_{B-L}$. See also Eq. (37).

[11] The lepton asymmetry is partially converted to the baryon asymmetry due to the ''sphaleron'' effects and the baryon asymmetry is given by $n_B/s = 8/23 |n_L/s|$ [26]. In this case, the evaporation of the associated Q balls (L balls) must be completed by the electroweak phase transition. This is possible, as we will see later, as long as $T_R \gtrsim 10^4$ GeV.

$$f_i|\phi|_0 < T, \quad c_i f_i T > H. \tag{22}$$

When the energy density of the universe is dominated by the oscillating inflaton, its decay produces the dilute plasma with $T \simeq (HT_R^2 M_*)^{1/4}$ [28]. This leads to the following sufficient condition to avoid the early oscillation of the AD field due to the thermal mass terms:

$$T_R < \frac{f_i}{c_i^{1/2}} M_* \left(\frac{|\phi_0|}{M_*} \right)^{3/2}. \tag{23}$$

This condition can be easily satisfied as long as $T_R \lesssim 10^7$ GeV, even if the cases where the coupling constants $f_i = \mathcal{O}(10^{-5})$ are present.

Second, let us investigate another thermal effect which is pointed out in Ref. [29]. The field, which has an effective mass $f_i|\phi| > T$, changes the trajectories of the running coupling constants of the light fields to which it couples. This effect produces the following potential for the AD field:

$$\delta V(\phi) = aT^4 \log\left(\frac{f_i^2 |\phi|^2}{T^2} \right)\Bigg|_{f_i|\phi|>T}, \tag{24}$$

where a is a constant which is given by the fourth power of gauge and/or Yukawa coupling constants, and it is at most $|a| = \mathcal{O}(10^{-2})$. Following the same method developed in Ref. [20], one can easily show that the above thermal effect does not become relevant if the following condition is satisfied:

$$T_R \lesssim \left(\frac{1}{|a|} \right)^{1/2} \left(\frac{m_\phi}{M_*} \right)^{1/2} |\phi_0|. \tag{25}$$

Hence, the early oscillation or trapping caused by the above thermal effect can also be avoided as long as $T_R \lesssim 10^7$ GeV. Thus, the estimation in the second line of Eq. (21) can be applied for a wide range of the reheating temperature $T_R \lesssim 10^7$ GeV in which we are free from the cosmological gravitino problem [24].

III. THE *Q*-BALL DECAY

In this section, we estimate the size of the associated Q balls and the conditions that the Q balls can evaporate or decay well above the freeze-out temperature of the LSP, which is crucial to avoid overclosing the universe. First, let us estimate the typical size of the associated Q balls following the methods in Refs. [10,12]. The relevant scalar potential for the AD field at the time of Q ball formation is given by

$$V(\phi) = m_\phi^2 \left[1 + K \log\left(\frac{|\phi|^2}{M_G^2} \right) \right] |\phi|^2, \tag{26}$$

where M_G is the renormalization scale at which m_ϕ is defined, and the K term represents the one-loop corrections dominantly from gaugino loops, and the value of K is estimated in the range from -0.01 to -0.1 [11,12,18]. The instability band for the AD field can be obtained from the equations for the linearized fluctuations and is given by [10]

$$0 < \frac{k^2}{R^2} < 3m_\phi^2 |K|, \tag{27}$$

where k is the comoving momentum of the fluctuations of the AD field. The best amplified mode is given by the center of the band,

$$\left(\frac{k^2}{R^2} \right) \max \simeq \frac{3}{2} m_\phi^2 |K|, \tag{28}$$

and it corresponds exactly to the Q-ball size which is estimated analytically using the Gaussian profile of the Q ball [12] as $R_Q \simeq \sqrt{2}/m_\phi |K|^{1/2}$. For the best amplified mode, the perturbations $\delta\phi$ grow according to the following equation:

$$\left| \frac{\delta\phi}{\phi} \right| \simeq \left| \frac{\delta\phi}{\phi} \right|_0 \exp\left(\int dt \frac{3}{4} m_\phi \left| K \right| \right), \tag{29}$$

and enter in nonlinear regimes when the Hubble parameter becomes $H_{non} = m_\phi |K|/2\alpha$, where $\alpha = \log(|\phi/\delta\phi|_0) \approx 30$. The Q balls are formed at this time, and hence the typical charge of a single Q ball can be estimated as

$$\begin{aligned} Q &\approx \frac{4}{3} \pi R_Q^3 \times n(t_{non}) \\ &\approx \frac{8\sqrt{2}\pi\beta}{27} \frac{|K|^{1/2}}{\alpha^2} |a_m| \delta_{eff} \left(\frac{m_{3/2}}{m_\phi} \right) \left(\frac{M_*}{m_\phi} \right)^3 \left(\frac{M_*}{M} \right)^3 \left(\frac{|S\phi^6|_0}{M_*^7} \right) \\ &\approx \frac{8\sqrt{2}\pi}{3\alpha^2} |K|^{1/2} B \left(\frac{M_*^2}{m_\phi T_R} \right) \\ &\approx 10^{16} \left(\frac{1 \text{ TeV}}{m_\phi} \right) \left(\frac{10^5 \text{ GeV}}{T_R} \right), \end{aligned} \tag{30}$$

where the subscript non denotes the time when the perturbations of the best amplified mode become nonlinear, and $B = n/s \sim 10^{-10}$, and we have used the fact that the Q balls are formed before the reheating process of inflation is completed. In the third line of Eq. (30), we have assumed that there are no additional entropy productions and have used Eq. (21) to connect the baryon or lepton number density when $H = H_{non}$ and the present baryon asymmetry. Note that the third line of Eq. (30) is also applicable for other Affleck-Dine baryogenesis scenarios as long as there are no additional entropy productions [12].

Now, we discuss the evaporation of the Q balls following the methods in Refs. [30,31]. Although the most part of a Q ball is decoupled from the thermal bath outside, the outer region of the Q ball is thermalized since particles in the surrounding plasma can penetrate into this region. The radius of the thermalized region inside the Q ball is estimated as

$$R_T = \gamma R_Q \equiv \left[\log\left(\frac{f_i|\phi(0)|}{T} \right) \right]^{1/2} R_Q. \tag{31}$$

Here, we have used the Gaussian profile of the Q ball; $|\phi(r)| \simeq |\phi(0)| \exp(-r^2/R_Q^2)$. Note that the evaporation of the baryonic or leptonic charge from the outer region of the Q ball is suppressed by diffusion effects as emphasized in Ref. [31] and is described by the following equation:

$$\Gamma_{\rm diff} = \frac{dQ}{dt} \simeq -4\pi D R_T \mu_Q(T) T^2, \tag{32}$$

where $D = \xi/T$ with $\xi \simeq 4-6$, and μ_Q is the chemical potential of the Q ball and it is in the range $m_\phi \lesssim \mu_Q(T) \lesssim T$. From Eqs. (32) and the relation between the cosmic time and temperature, one can easily show that the evaporation is more efficient for lower temperatures and the total amount of the evaporated charge is mainly provided at $T \simeq m_\phi$, since for temperatures below this value, the evaporation is exponentially suppressed by the Boltzmann factor $\exp(-m_\phi/T)$. Thus, the total amount of the evaporated charge from a single Q ball can be estimated as

$$\triangle Q \approx \frac{4\sqrt{2}\pi}{\zeta^{1/2}|K|^{1/2}} \xi \gamma \left(\frac{M_*}{m_\phi}\right) \approx 10^{18} \left(\frac{0.01}{|K|}\right)^{1/2} \left(\frac{1\ \text{TeV}}{m_\phi}\right), \tag{33}$$

where $\zeta = \pi^2 g_*/90$. Here, we have assumed $T_R \gtrsim m_\phi$. If the initial charge of a Q ball is larger than this value, the remaining charge is emitted by the decay into light fermions. The upper bound on the decay rate into light fermions is given by [32]

$$\left|\frac{dQ}{dt}\right|_{\rm fermion} \leqslant \frac{\omega^3 \mathcal{A}}{192\pi^2}, \tag{34}$$

where $\mathcal{A}$ is the surface area of the Q ball, and $\omega \simeq m_\phi$. The upper bound is likely to be saturated for Q balls with $\phi(0)$ much larger than m_ϕ, which has been found from numerical calculations [32]. However, there might be an enhancement factor f_s due to the decay into lighter scalars, which is expected to be at most $\mathcal{O}(10^3)$. Thus, the decay rate of a Q ball can be written as [12]

$$\frac{dQ}{dt} = f_s \left(\frac{dQ}{dt}\right)_{\rm fermion}. \tag{35}$$

By integrating Eq. (35), we can obtain the decay temperature of a Q ball in the following form:

$$T_d \lesssim 2\ \ \text{GeV} \left(\frac{0.01}{|K|}\right)^{1/2} \sqrt{f_s} \left(\frac{m_\phi}{1\ \ \text{TeV}}\right)^{1/2} \left(\frac{10^{20}}{Q}\right)^{1/2}. \tag{36}$$

If the decay temperature of a Q ball is well above the freeze-out temperature of the LSP, $T_f \approx m_\chi/20$, the annihilation of the produced LSP's effectively takes place. We see from Eqs. (30), (33), and (36) that if the reheating temperature of inflation satisfies $T_R \gtrsim 10^3$ GeV, the Q ball charge is estimated as $Q \lesssim 10^{18}$, and hence we are free from the LSP over-production problem. From Eq. (21), one can see that this condition can be easily satisfied in our scenario with the interesting value of the $U(1)_{B-L}$ breaking scale $v \simeq (2-7)\times 10^{14}$ GeV. Note that the breaking scale of the $U(1)_{B-L}$ can be determined almost independently of the reheating temperature of the inflation as long as 10^3 GeV $\lesssim T_R \lesssim 10^7$ GeV.

IV. DISCUSSION AND CONCLUSIONS

In this paper, we perform an analysis of Affleck-Dine baryogenesis/leptogenesis scenarios in the context of the minimal extension of the MSSM in which the $U(1)_{B-L}$ symmetry is gauged. We find that all of the relevant flat directions can be lifted at the breaking scale of the $U(1)_{B-L}$ by the $U(1)_{B-L}$ D-term contribution, and the LSP over-production problem can easily be avoided. The required baryon asymmetry from the big-bang nucleosynthesis can be obtained in a wide range of the reheating temperature of inflation, 10^3 GeV $\lesssim T_R \lesssim 10^7$ GeV, which are low enough to avoid the cosmological gravitino problem [24]. Surprisingly enough, although the breaking scale of the $U(1)_{B-L}$ symmetry is determined totally independently of the right-handed neutrino masses, the obtained scale from the baryon asymmetry is quite consistent with the scale suggested from the seesaw mechanism to explain the recent neutrino-oscillation experiments [25].

We also comment here on other possibilities to avoid the LSP over-production problem in Affleck-Dine baryogenesis/leptogenesis scenarios which use dimension six nonrenormalizable operators in the superpotential. We make the associated Q balls sufficiently small if we can suppress the baryon asymmetry in the initial stage of the AD field oscillation (which is the relevant epoch for the formation of Q balls). We find that it is, in fact, possible if the effective cutoff scales are lower than the order of the 10^{16-17} GeV and the reheating temperature is in the range 10^3 GeV $\lesssim T_R \lesssim 10^5$ GeV.[12] This may be an interesting possibility since such a relatively low cutoff scale is suggested in the $\mathcal{M}$

[12]These parameter regions for the effective cutoff scale and the reheating temperature can be understood as follows. Without a gauged $U(1)_{B-L}$, the evolution of the AD field before it starts the coherent oscillation is given by $|\phi| \simeq (HM^3)^{1/4}$, which is determined from the negative Hubble-mass term and the dimension ten operators. Then, the baryon asymmetry can be calculated following a similar method in Sec. II as

$$\frac{n_B}{s} \simeq \frac{2\beta}{9} \delta_{eff} \frac{T_R m_{3/2} |a_m|}{M_*^2} \left(\frac{M}{m_\phi}\right)^{3/2} \sim 10^{-11} \left(\frac{T_R}{10\ \ \text{GeV}}\right) \left(\frac{m_{3/2}}{1\ \ \text{TeV}}\right) \times \left(\frac{1\ \ \text{TeV}}{m_\phi}\right)^{3/2} \left(\frac{M}{M_*}\right)^{3/2}, \tag{37}$$

where the definitions are the same with those in Sec. II. From Eqs. (30), (33), and (37), we can see $M \lesssim 10^{16-17}$ GeV is required for the evaporation of the associated Q balls. On the other hand, there might appear strong thermal effects for $M \lesssim 10^{15}$ GeV, and the above estimation might be drastically changed.

theory [33]. Other possibilities such as $f_B \ll 1$ or to use the NMSSM with a light singlino are investigated in Ref. [21].

Finally, we point out that our scenario can also work in the case of anomaly mediated SUSY breaking [3,4]. Note that there is a general serious problem [34] for Affleck-Dine baryogenesis and leptogenesis scenarios in the case of anomaly-mediated SUSY breaking. Without a gauged $U(1)_{B-L}$ symmetry, the relevant scaler potential for the AD field which is lifted by n-dimensional superpotential is given by

$$V=(m_\phi^2-cH^2)|\phi|^2-\frac{m_{3/2}}{nM^{n-3}}(a_m\phi^n+\text{H.c})+\frac{1}{M^{2(n-3)}}|\phi|^{2(n-1)}, \tag{38}$$

where the gravitino mass $m_{3/2}$ is much larger than the soft mass for the AD field m_ϕ in the anomaly mediation models. Because of the presence of the large A terms, there appears a global minimum displaced from the origin, $|\phi|_{\min} \sim (m_{3/2}M^{n-3})^{1/n-2}$, and the AD field is expected to be trapped in this minimum during its slow rolling regime. In our scenario, the scalar potential is given by Eq. (13), but with much larger gravitino mass. The field value of the top of the hill in front of this global minimum is given by

$$|\phi|_{\text{hill}} \sim \left[m_\phi M^3 \frac{m_\phi}{m_{3/2}} \left(\frac{M_*}{\langle |S| \rangle} \right) \right]^{1/4}. \tag{39}$$

Thus, as far as the initial amplitude of the AD field is smaller than this value $|\phi|_{\text{hill}}$, the Affleck-Dine mechanism can work without any difficulty, and in fact, this is what happens in our scenario with the gauged $U(1)_{B-L}$.

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APPENDIX: THE Q BALLS IN THE AFFLECK-DINE BARYOGENSIS AND LEPTOGENESIS SCENARIOS WITH $W=0$

In this appendix, we will briefly investigate the features of the Q balls formed in the Affleck-Dine baryogenesis and leptogenesis scenarios without superpotential. In this case, the AD field develops its field value as large as the reduced Planck scale during the inflation and starts its coherent oscillation when the Hubble parameter becomes smaller than its soft mass. There are no relevant thermal effects because the masses of the fields which couple to the AD field are much larger than the temperature. The baryon and/or lepton number density when the AD field starts its coherent oscillation is given by

$$n(t_0) \approx \beta \delta_{eff} \left(\frac{m_{3/2}}{m_\phi} \right)^2 m_\phi M_*^2, \tag{A1}$$

where $\delta_{eff}=\mathcal{O}(1)$. Based on a similar argument which leads to Eq. (30), the typical size of a single Q ball is estimated as

$$Q \approx \frac{4}{3}\pi R_Q^3 \left(\frac{H_{non}}{m_\phi} \right)^N \times n(t_0), \tag{A2}$$

where $N=3/2$ for $T_R > \zeta^{-1/4}\sqrt{m_\phi M_*}$ and $N=2$ for $T_R < \zeta^{-1/4}\sqrt{m_\phi M_*}$. Even if we assume the Q ball is formed in the inflaton-dominated epoch, i.e., $N=2$, the typical size of charge is $Q \sim 10^{26}$. From Eqs. (33) and (36), it is clear that these Q balls cannot evaporate, and its decay temperature is well below the freeze-out temperature of the LSP.

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