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Affleck-Dine baryogenesis after thermal brane inflation

Tomohiro Matsuda*

Laboratory of Physics, Saitama Institute of Technology, Fusaiji, Okabe-machi, Saitama 369-0293, Japan

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We propose a new scenario of Affleck-Dine baryogenesis in the context of theories with large extra dimensions. In this paper we consider baryogenesis after thermal brane inflation and show how our mechanism works. We specifically consider models in which supersymmetry is broken at the distant brane.

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I. INTRODUCTION

The production of net baryon asymmetry requires baryon number violating interactions, C and CP violation and a departure from thermal equilibrium [1]. The first two of these ingredients are naturally contained in grand unified theories (GUTs) or other string-motivated scenarios, and the third can be realized in an expanding universe where it is not uncommon that interactions come in and out of equilibrium, producing the stable heavy particles or cosmological defects. In the original and simplest model of baryogenesis [2], a heavy GUT gauge or Higgs boson decays out of equilibrium producing a net baryon asymmetry.

Another mechanism for generating the cosmological baryon asymmetry in supersymmetric theories is proposed by Affleck and Dine [3] who utilized the decay of the scalar condensate along the flat direction. This mechanism is a natural product of supersymmetry, which contains many flat directions that break $U(1)_B$. The scalar potential along this direction vanishes identically when supersymmetry breaking is not induced. Supersymmetry breaking lifts this degeneracy,

$$V \simeq m_{soft}^2 |\phi|^2 \tag{1.1}$$

where m_{soft}^2 is the supersymmetry-breaking scale and ϕ is the direction in the field space corresponding to the flat direction. For the large initial value of ϕ, a large baryon number asymmetry may be generated if the condensate of the field breaks $U(1)_B$. The mechanism also requires the presence of baryon number-violating operators that may appear through higher dimensional A terms. The decay of these condensates through such an operator can lead to a net baryon asymmetry. In the most naive consideration the baryon asymmetry is computed by tracking the evolution of the sfermion condensate in the flat direction of the supersymmetric standard model. Considering a toy model with the potential

$$V(\phi,\phi^\dagger) = m_{soft}^2 |\phi|^2 + \frac{1}{4}[\lambda \phi^4 + \text{H.c.}], \tag{1.2}$$

the equation of motion becomes

$$\ddot{\phi} + 3H\dot{\phi} = -m_{soft}^2 \phi + \lambda (\phi^\dagger)^3. \tag{1.3}$$

The baryon (or lepton) number density is given by

$$n_B = q_B(\dot{\phi}^\dagger \phi - \phi^\dagger \dot{\phi}), \tag{1.4}$$

where q_B is the baryon (or lepton) charge carried by the field. Now one can write down the equation for the baryon number density

$$\dot{n}_B + 3Hn_B = 2q_B Im[\lambda(\phi^\dagger)^4]. \tag{1.5}$$

Integrating this equation, one can obtain the baryon (or lepton) number produced by the Affleck-Dine oscillation. For a large initial amplitude, the produced baryon number is estimated as $n_B \simeq (4q_B|\lambda|/9H)|\phi_{ini}|^4 \delta_{eff}$, where δ_{eff} is the effective CP violation phase of the initial condensate. This crude estimation suggests that by generating some angular motion one can generate a net baryon density.

In the conventional scenario of Affleck-Dine baryogenesis, one should assume large $H > m_{3/2}$ before the time of Affleck-Dine baryogenesis so that the flat directions are destabilized to obtain the large initial amplitude of baryon-charged directions.

Although it seems plausible that Affleck-Dine baryogenesis generates the baryon asymmetry of the Universe, there are some difficulties in the naive scenario. The formation of a Q ball [4] is perhaps the most serious obstacle that puts a serious constraint on the baryon number density at the time of Q-ball formation. Q balls are formed due to the spatial instability of the Affleck-Dine field, and have been shown by numerical calculations that they absorb almost all the baryonic charges in the Universe when they form [5,6]. This means that the baryon asymmetry of the Universe in the later period must be provided by decaying Q balls. In general, the stability of Q balls are determined by their charge that are inevitably fixed by Affleck-Dine mechanism itself. The reason is that the formation of Q balls occurs almost immediately, which makes it hard to expect any additional diluting mechanism before Q-ball formation. The point is that in general Affleck-Dine baryogenesis the initial baryon number density becomes so huge that the produced Q balls become stable. The stable Q balls that produce the present baryon asymmetry of the Universe by their decay are dangerous, because such Q balls can also produce dangerous relics at the same time when they decay to produce the baryons. The decay temperature of the associated huge Q balls becomes in

*Email address: matsuda@sit.ac.jp

general much lower than the freeze-out temperature of the dangerous lightest supersymmetric particle, which causes serious constraint.

Another obstacle is the problem of the early oscillation caused by the thermalization [7]. When the fields that couple to the Affleck-Dine field are thermalized, they induce the thermal mass to the Affleck-Dine field. The early oscillation starts when the thermal mass term exceeds the destabilizing mass. The serious constraint appears because the destabilizing mass, which is about the same order of the Hubble parameter, is in general much smaller than the temperature of the plasma.

However, in our model these difficulties do not appear since the mechanism of the destabilization of the Affleck-Dine field is not a consequence of large Hubble parameter. The size of the Q ball is naturally suppressed, since our mechanism does not produce a huge baryon number density after Affleck-Dine baryogenesis.

What we will consider in this paper is a mechanism in which Affleck-Dine mechanism is realized after thermal brane inflation [8]. Before discussing the baryogenesis with extra dimensions, we must first specify the scenario of the early Universe to a certain extent. In this paper we consider Affleck-Dine baryogenesis after thermal brane inflation.[1] We show how our mechanism works in models with supersymmetry breaking at the distant brane. Here we consider two different cases for supersymmetry breaking. In one case we assume an alternative source of supersymmetry breaking on the distant brane, and in the other case we deal with the realistic bulk field mediation of supersymmetry breaking. It is often the case that the brane distance is used to control the direct contact terms that produce unwanted soft terms preventing the flavor changing neutral current (FCNC) bound. Our mechanism is expected to work in these models if the relevant brane distance is reduced during a period after inflation. In addition to the thermal inflation model that we have considered in this paper, there are many models in which the temporally reduced extra dimension is used to prevent difficulties related to the large extra dimensions [10].

II. THERMAL BRANE INFLATION

In this section we make a brief review of thermal brane inflation proposed by Dvali [8]. The following conditions are required so that the mechanism functions.

(1) Exchange of the bulk modes such as graviton, dilaton or Ramond-Ramond (RR) fields governs the brane interaction at the large distance.

(2) In the case when branes initially come close, bulk modes are in equilibrium and their contribution to the free energy can create a positive T^2 mass term for ϕ to stabilize the branes on top of each other until the Universe cools down to a certain critical temperature[2] $T_c \sim m_s$. Here m_s represents the negative curvature of ϕ at the origin determined by the supersymmetry breaking, and ϕ is the moduli field for the brane distance.

The resulting scenario of thermal inflation is straightforward. Assuming that there was a period of an early inflation with a reheat temperature $T_R \sim M$, and at the end of inflation some of the repelling branes sit on top of each other stabilized by the thermal effects, one can obtain the number of e-foldings

$$N_e = \ln\left(\frac{T_R}{T_c}\right). \qquad (2.1)$$

Taking $T_R \sim 10$ TeV and $T_c \sim 10^3 - 10$ MeV, one finds $N_e \sim 10$–15, which is consistent with the original thermal inflation [11] and is enough to get rid of unwanted relics.[3]

III. AFFLECK-DINE BARYOGENESIS AFTER THERMAL BRANE INFLATION

In this section we show how to realize Affleck-Dine (AD) baryogenesis after thermal brane inflation. Our model requires the mechanism of supersymmetry breaking at the distant brane. To proceed, we should first discuss the mechanism of supersymmetry breaking. In our model the negative soft term is not a simple consequence of the large Hubble parameter, but rather related to the distance between the matter brane and the supersymmetry-breaking brane. We should also discuss the origin of the baryon number violating A terms, which plays the crucial role in Affleck-Dine baryogenesis. Because of the constraint from proton stability, an additional mechanism for suppressing dangerous higher dimensional A terms is always required when the fundamental scale is much lower than the Planck scale.

In the oldest version of supergravity mediation, it is assumed that all higher-dimension operators that directly connect the fields in the hidden sector with the ones in the observable sector are present but suppressed only by powers of $1/M_4$, where M_4 denotes the Planck mass in four dimensions. In this case the required soft supersymmetry-breaking

[1] In theories with extra dimensions there are two possible choices for the Affleck-Dine field. It could either be a brane field localized on a brane or a bulk field. In Ref. [9], it is discussed that naive realization of Affleck-Dine mechanism with a brane field cannot produce sufficient baryon number.

[2] The authors of Ref. [8] considered open string modes stretched between the different branes. If the branes are on top of each other, these string modes that get mass when the brane distance grows are in equilibrium and their contribution to the free energy creates a positive T^2 mass term so that the resulting curvature becomes positive.

[3] In this paper we also consider situations where the reheating temperature after the first inflation is as high as $T_R \sim 10^{10}$ GeV, and the critical temperature $T_c \sim 10^2$ GeV. Unlike the original model of thermal brane inflation, large extra dimensions are not specially supposed in this paper. Since we are taking interest in whether our mechanism of baryogenesis works, we also deal with the case where the thermal brane inflation itself is not a necessary ingredient to solve the cosmological problems. In such cases, the concern should be whether there can be a short period of thermal brane inflation that enables our mechanism of Affleck-Dine baryogenesis to work.

term is given by the higher-dimension terms of the form

$$L_{soft} \sim \int d^4\theta \frac{1}{M_4^2} X^\dagger X Q^\dagger Q. \tag{3.1}$$

Here X is a chiral superfield in the hidden sector whose F component F_X breaks supersymmetry. Q is a matter field in the visible sector. Higher-dimensional operators in the superpotential $W_A \sim (1/M_p^{n+3})\Phi^{n+3}$ produce the A terms and determines the phase of the AD direction at large $\langle\Phi\rangle$:

$$L_A \sim \int d^4\theta \left(\frac{1}{M_4^{n+3}} X^\dagger X \Phi^{n+3} + \text{H.c.} \right) + \int d^2\theta \left(\frac{1}{M_4^{n+1}} X \Phi^{n+3} + \text{H.c.} \right), \tag{3.2}$$

where $n \geqslant 1$ and Φ represents the flat direction.

Contact terms of the similar form appear in the models of extra dimensions, where M_4 is replaced by the fundamental scale M that is much lower than M_4. On the other hand, the contact terms connecting the fields in the hidden and the observable branes are suppressed because they are localized along the extra dimension. In these cases the supersymmetry breaking is mediated by bulk fields such as scalar fields [12,13] or fermions [14] where the scale of the supersymmetry breaking in the hidden brane can be as large as the fundamental scale of the higher-dimensional theory, while the direct soft terms for the standard model sfermions are suppressed.

A. Simplest toy model

For the simplest toy model we consider an example where the fundamental scale M is as low as 10 TeV and the realistic supersymmetry breaking is realized within the matter brane without specific fine tuning. In addition to these simplest settings, we also include a distant brane where the supersymmetry is maximally broken by an auxiliary component of a localized field $|F_X|^{1/2} \sim M$. In such a case the effect of F_X on the matter brane is expected to be exponentially suppressed because they are localized at the distant brane. The soft terms are given by

$$V(\phi_{AD}) \sim \left[m_{soft}^2 + c \left(\frac{|F_X|}{M} \right)^2 e^{-M r_{susy}} \right] |\phi_{AD}|^2. \tag{3.3}$$

Here ϕ_{AD} is the flat direction of the Affleck-Dine mechanism, and r_{susy} is the distance between the matter brane and the hidden supersymmetry-breaking brane on which F_X is localized. m_{soft} denotes the supersymmetry-breaking induced on the matter brane, which is assumed to be a constant. When two branes sit on their true positions, the second term is negligible. On the other hand, when the hidden brane stays on top of the matter brane during thermal brane inflation, then the supersymmetry breaking of order F_X/M is induced on the matter brane by the direct contact terms. Assuming that the effective soft mass appears with the negative sign (i.e., $c<0$), the flat direction ϕ_{AD} is destabilized during thermal inflation if $m_{soft} < |F_X|/M$. At the same time A terms are modified to generate the required misalignment of the phase. Here we assume that the A term is effectively given by using the four-dimensional Planck mass,

$$V_A \simeq \left(\frac{a_0 m_{soft}}{M_p} + \frac{a_1 |F_X| e^{-M r_{susy}}/M}{M_p} \right) \phi_{AD}^4 \tag{3.4}$$

where a_0 and a_1 are constants of $O(1)$. The situation here is very similar to the original Affleck-Dine baryogenesis. The sole difference is that the supersymmetry is not induced by the Hubble parameter, but is induced by the brane distance. The resultant baryon to entropy ratio is[4]

$$\frac{n_B}{s} \sim \frac{T_{R2}}{M_p H_o \rho_I} |a m_{soft} (\phi_{AD}^i)^4| \delta_{eff} \tag{3.5}$$

where T_{R2} is the reheating temperature after thermal brane inflation, and ϕ_{AD}^i is the initial amplitude of ϕ_{AD}. H_o denotes the Hubble parameter when the AD oscillation starts, which can be taken to be $H_o \leqslant H_I = M^2/M_p$. It is naturally assumed that the initial amplitude is $\phi_{AD}^{ini} \sim M$, and the inflaton density is still $\rho_I \sim M^4$ at the beginning of the oscillation. Then we obtain

$$\frac{n_B}{s} \sim 10^{-10} \left(\frac{T_{R2}}{10 \ \text{MeV}} \right) \left(\frac{10^{-8} \ \text{GeV}}{H_o} \right) \tag{3.6}$$

which is the most naive result, but is enough to explain the origin of the baryon asymmetry of the present Universe.[5]

B. Realistic mediation of supersymmetry breaking

Here we consider gaugino mediation as a more realistic example of such ''hidden'' supersymmetry breaking.[6] In this case the minimal supersymmetric standard model (MSSM) scalar mass squareds derived from five-dimensional Feynman diagrams are suppressed relative to the gaugino masses by at least a loop factor when the brane distance is larger than M^{-1}.

Even in the limit of small brane distance, it does not exceed the masses generated from the four-dimensional renormalization-group evolution between the compactification scale and the weak scale.[7] This conclusion is generic and also holds for the other soft parameters such as A terms.

Besides the contributions from five-dimensional Feynman diagrams of gauginos propagating through extra dimension, there are direct contact terms that can destabilize the flat direction during thermal brane inflation.

[4]See Ref. [15] for more detail.

[5]Modifications of parameters are allowed, but in general they are strongly model dependent. The magnitude of the A term can be modified at the time of AD oscillation, which we shall discuss in the next paragraph.

[6]Here we have assumed that the gaugino can propagate only one extra dimension that is about $10-10^2$ times larger than M^{-1} [14].

[7]Details of the calculations are given in Ref. [17].

Here we consider two sources of supersymmetry breaking, the four-dimensional effect and the direct contact term. Assuming that the soft terms of the relevant flat direction is produced by these two sources, it takes the following form [16]:

$$V(\phi_{AD})\sim\left[c_1\left(\frac{g_4^2}{(4\pi)^2}\right)^2\left(\frac{|F_X|}{M}\right)^2 + c_3\left(\frac{|F_X|}{M}\right)^2 e^{-Mr_{susy}}\right]|\phi_{AD}|^2 \tag{3.7}$$

for small ϕ_{AD} and

$$V(\phi_{AD})\sim c_2\left(\frac{g_4^2}{(4\pi)^2}\right)^2(|F_X|)^2\left(\ln\frac{|\phi_{AD}|^2}{M^2}\right)^2 + c_3\left(\frac{|F_X|}{M}\right)^2 e^{-Mr_{susy}}|\phi_{AD}|^2 \tag{3.8}$$

for large ϕ_{AD}. Here ϕ_{AD} is the flat direction of Affleck-Dine mechanism, and M is the fundamental scale. r_{susy} is the distance between the supersymmetry-breaking brane and the matter brane.

If the supersymmetry-breaking hidden brane stays on top of our brane during thermal brane inflation, the supersymmetry breaking on our brane during this period is naturally the order of $O(|F_X|/M)$ because the $e^{-Mr_{susy}}$ factor in the direct contact term is $O(1)$. Relevant soft mass is then given by

$$m^2(\phi_{AD})\sim c_1\left(\frac{g^2}{(4\pi)^2}\right)^2\left(\frac{|F_X|}{M}\right)^2 + c_3\left(\frac{|F_X|}{M}\right)^2 e^{-Mr_{susy}} \tag{3.9}$$

for $\phi_{AD}<M$. It is obvious that the source of supersymmetry breaking during thermal brane inflation is in general different from the one at the true vacuum. At this time the flat directions on the observable brane are lifted or destabilized by the supersymmetry breaking induced by the direct contact terms, which will soon disappear as soon as the brane distance grows. Then the situation is similar to the conventional Affleck-Dine baryogenesis, where the destabilization is induced by the alternative supersymmetry breaking induced by the inflaton. Assuming that the direct supersymmetry breaking destabilizes the flat direction with the negative soft mass, which corresponds to taking the constant $c_3<0$, the potential of the Affleck-Dine flat direction at the end of thermal inflation is given by

$$V_{soft}(\phi_{AD})\sim -|c_3|\left(\frac{|F_X|}{M}\right)^2|\phi_{AD}|^2. \tag{3.10}$$

This negative soft mass disappears soon after the end of thermal brane inflation, as the brane distance r_{susy} grows.

The direct contact terms decreases exponentially, while terms produced by the four-dimensional effect are not modified by r_{susy}, because the supersymmetry-breaking gaugino mass is determined by the size of the relevant extra dimension that is assumed to be a constant during thermal brane inflation. Then there should be an oscillation of the Affleck-Dine field that starts at $r_{susy}\sim M^{-1}$.

We should also consider another important ingredient of Affleck-Dine baryogenesis, the evolution of the A term. To discuss the magnitude of A terms during thermal brane inflation, we must first discuss a concrete model for suppressing the baryon number-violating interactions. The most naive idea is to assume that the baryon number is maximally broken on the hidden brane and its effect on our brane is exponentially suppressed by e^{-r_B} where r_B is the distance between the hidden brane and the matter brane [18–20]. A popular mechanism for explaining the smallness of the observed Yukawa couplings or baryon number violating operators is to expect higher dimension operators of the generic form

$$\mathcal{O}\sim\lambda\left(\frac{\chi}{M}\right)^k\mathcal{O}_{MSSM} \tag{3.11}$$

with $\lambda\sim O(1)$. If $\epsilon\sim\chi/M$ is small, the small couplings in these operators are understood as the small parameter ϵ. The smallness of ϵ is understood if the shined value of $\langle\chi\rangle$ is assumed on matter brane [20]. Assuming that $\langle\chi\rangle\sim M$ at the distant brane and their mass in the bulk is about $\sim M$, the suppression factor is given by the shining method[8]

$$\epsilon\sim\frac{e^{-Mr_B}}{r_B^{n_E-2}} \tag{3.12}$$

for $n_E>2$ and $r_BM\gg1$, where n_E denotes the number of the relevant extra dimensions. For $n_E=2$ and $r_BM\gg1$,

$$\epsilon\sim\frac{e^{-Mr_B}}{\sqrt{Mr_B}}. \tag{3.13}$$

Thus one can obtain the e^{-Mr_B} suppression for each ϵ. Here r_B denotes the distance between the baryon number-breaking brane and the matter brane. In this case, because of the suppression ϵ^k, baryon number-violating A terms are safely suppressed by the exponential factor at the true vacuum in order not to produce dangerous operators. On the other hand, because the suppression factor originates from the brane distance r_B, such A terms are not suppressed when branes are on top of each other.[9] Let us consider an example where a higher dimensional term with the lowest k determines the phase of the Affleck-Dine condensate in the true vacuum, while other terms dominate when $r_B=0$. The phases of these direct contacting A terms are in general different from the one at the true vacuum, thus producing the misalignment of the phase during thermal brane inflation. Because these alternative contributions become tiny right after the end of ther-

[8]See Ref. [20] for more detail.

[9]Of course one can assume that the baryon number-breaking hidden brane is identical to the supersymmetry-breaking hidden brane. In such a case, the brane distance r_{susy} is identified with r_B.

mal brane inflation, misalignment of the phase is expected to appear just after thermal brane inflation.

Of course one can expect the case where the thermal brane inflation does not modify the baryon number-violating operators in the superpotential. This happens when the smallness of the operator is produced by other mechanisms that are not relevant to the brane distance, or in the case where r_B is not modified during thermal inflation. In this case the modification of the A term is induced only by the supersymmetry breaking, which is precisely the same as what happens in the conventional Affleck-Dine baryogenesis.

In both cases, one can expect that the baryogenesis starts at $r_B \sim M^{-1}$ or $r_{susy} \sim M^{-1}$, where the exponential suppression becomes significant. The calculation of the resultant baryon to entropy ratio is similar to the conventional Affleck-Dine baryogenesis. Here we assume that the A term of the form

$$V_A \simeq \frac{a m_{soft}}{M} \phi_{AD}^4 \tag{3.14}$$

is already recovered at the beginning of the AD oscillation. Taking the initial amplitude $\phi_{AD}^{ini} \sim M$, we obtain [15]

$$\frac{n_B}{s} \sim \frac{T_{R2}}{H_o} \frac{|a m_{soft} (\phi_{AD}^i)^4|}{M \, \rho_I} \delta_{eff}. \tag{3.15}$$

Here the inflaton density is denoted by $\rho_I \sim M^4$. Then we can obtain

$$\frac{n_B}{s} \sim 10^{-10} \left(\frac{a}{10^{-7}} \right) \left(\frac{T_{R2}}{10 \text{ GeV}} \right) \left(\frac{10^8 \text{ GeV}}{M} \right)^3. \tag{3.16}$$

Of course in some cases the inflaton density may be determined by the scale of the supersymmetry-breaking auxiliary component F_X, such as $\rho_I \sim |F_X|^2$. In this case the baryon to entropy ratio becomes about $O(|M/F_X^{1/2}|^4)$ times larger than Eq. (3.16).

The most significant difference from the conventional Affleck-Dine baryogenesis with extra dimensions is the absence of the problematic suppression factor that makes it impossible to realize Affleck-Dine baryogenesis on the brane [9].

IV. CONCLUSIONS AND DISCUSSIONS

In this paper we have considered an alternative mechanism of Affleck-Dine baryogenesis that starts after thermal brane inflation. Our mechanism works in models with supersymmetry breaking at the distant brane. The brane distance is required to be modified during thermal brane inflation in order to activate the alternative source of supersymmetry breaking. In addition to the thermal inflation that we have considered in this paper, there are many models in which the initially reduced extra dimensions are used to prevent difficulties related to the large extra dimensions [10]. Extensions to these models will be discussed in the next publication [21].

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