Affleck-Dine baryogenesis in the local domain

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Abstract

For Affleck-Dine baryogenesis to proceed, there must have been two types of phase transitions. One is the destabilized-stabilized phase transition of the flat direction, which is in general induced by the Hubble parameter. The other is the phase transition related to the A-term, which induces the misalignment of the relative phase of the flat direction. In the conventional Affleck-Dine baryogenesis they are supposed to start almost simultaneously. Of course these phase transitions can take place separately, but the latter must not be later than the former because the phase transition of the A-term can not produce any baryon number when there is no condensate of the relative charge.

In this paper we try to construct models where the original idea of Affleck-Dine baryogenesis is realized in a different way. We show examples in which the local domain of the false vacuum with the required condensate is formed after inflation and collapses in a safe way so that the domain wall problem is avoided. We also show examples where the phase transition of the A-term starts before the decay of the condensate. As in the conventional Affleck-Dine mechanism, the phase transition of the A-term produces baryon number in the local domain of the condensate. We construct scenarios where our mechanism produces sufficient baryon asymmetry of the Universe.

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1 Introduction

The production of net baryon asymmetry requires baryon number violating interactions, C and CP violation and a departure from thermal equilibrium[1]. In the original and simplest model of baryogenesis[2] a GUT gauge or Higgs boson decays out of equilibrium producing a net baryon asymmetry. Another mechanism of generating the cosmological baryon asymmetry in supersymmetric theories is proposed by Affleck and Dine[3] who utilized the decay of 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,

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

where $m_{soft}$ is the supersymmetry-breaking scale and $\phi$ is the direction in the field space corresponding to the flat direction. For large initial values of $\phi$, 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. Considering a toy model with the potential

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

the equation of motion becomes

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

The baryon (or lepton) number density is given by:

$$ n_B = q_B \left( \dot{\phi}^\dagger \phi - \phi^\dagger \dot{\phi} \right), $$

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

$$ \dot{n}_B + 3H n_B = 2q_B \text{Im} \left[ \lambda (\phi^\dagger)^4 \right]. $$
Integrating this equation, one can obtain the baryon (or lepton) number produced by the Affleck-Dine oscillation. For large initial amplitude, the produced baryon number is estimated as \( n_B \approx \frac{4q_{\lambda}^2}{H} |\phi_{ini}|^4 \delta_{eff} \), where \( \delta_{eff} \) 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 conventional mechanism of Affleck-Dine baryogenesis, one should assume large \( H > m_{3/2} \) before the time of Affleck-Dine baryogenesis to destabilize the flat directions and obtain the large initial amplitude of baryon-charged directions.

Although it seems plausible that Affleck-Dine baryogenesis generates baryon asymmetry of the Universe, there are some difficulties in the naive scenario. The formation of Q-ball[4] is perhaps the most serious obstacles 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-ball is determined by their charge, which is inevitably fixed by Affleck-Dine mechanism itself. The reason is that the formation of Q-ball 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 are 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 Q-balls becomes in general much lower than the freeze-out temperature of the dangerous lightest supersymmetric particle, which causes serious constraint.

Another obstacle is the problem of early oscillation caused by the thermalization[7]. When the fields that couple to the Affleck-Dine field is 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.
Now let us explain what we have in mind. In general Affleck-Dine baryogenesis the following conditions are required to produce net baryon asymmetry.

- There exists the field condensate $\phi_{AD} \neq 0$ charged with the baryon number. The condensate must decay during or after the baryon number production.

- The A-term that determines the initial phase of the condensate must be different from the one in the true vacuum. There must be the phase transition of the A-term that rotates the phase of the condensate. The rotation produces net baryon number.

In conventional models of Affleck-Dine baryogenesis it is assumed that the phase transition of the collapsing condensate and the phase transition of the A-term take place almost simultaneously during the inflaton-dominated era of the Universe. Both of them were supposed to be induced by the supersymmetry breaking of the decreasing Hubble parameter. In the scenario of the conventional Affleck-Dine baryogenesis, the A-term induced by non-zero $H$ displaces the CP-violating phase. As $H$ becomes smaller, other sources for the A-term dominate and the shift of the phase turns into the driving force in the angular direction, which produces net baryon number.

However, in some cases, these phase transitions do not start simultaneously. The most serious example is the early oscillation induced by the thermal mass of the flat direction, which starts before the phase transition of the A-term\(^7\). In such cases, the required baryon number is not produced. On the other hand, one can remedy this situation if the phase transition of the A-term starts earlier than the collapse of the condensate\(^7\).

Here we try to construct models where the required phase transitions are completely separated from the inflation. When one discusses the cosmological domain wall, one must take into account the large (but local) domains of the quasi-degenerated vacua. In generic situations the radius of the domain immediately becomes as large as the Hubble radius. If the required condensate exists in one of these domains, the first requirement $\phi_{AD} \neq 0$ is satisfied locally in the domain. Then our task is to construct a model where the A-term changes its phase before the collapse of the quasi-degenerated domain of the condensate.\(^2\)

\(^2\)The criteria for the cosmological domain wall is discussed in ref.\(^8\). If they ever existed, they must have disappeared at $t < t_c \sim (G\sigma)^{-1}\(^3\)$.\(^4\)
The most significant difference from the conventional Affleck-Dine baryogenesis is that the production of the baryon number starts at much later period of the evolution of the Universe.

In this paper we construct a concrete model in which the required condensate is realized in the domain of the quasi-degenerated vacua. Our model starts with the situation where the Universe is occupied by the domains of the false and the true vacua. We also construct models in which the phase transition of the A-term is realized at low energy scale. Let us consider a toy model with a potential in the domain of the false vacuum,

$$V(\phi, \phi^\dagger) = -m^2 |\phi_{\text{AD}}|^2 + \left[ \frac{a_o}{4M_p} \phi_{\text{AD}}^4 + h.c. \right]$$ (1.6)

where the vacuum expectation value of the Affleck-Dine field $\phi_{\text{AD}}$ is non-zero. $a_o$ represents the A-term in the true vacuum.\footnote{Of course one may consider the case where the destabilization of the field (1.6) appears inside cosmological defects such as walls and strings. Then there will be the volume factor suppression. The volume suppression is $\epsilon_{\text{vol}} \sim \Delta_{\text{AD}}/H^{-1}$ for domain walls and $\epsilon_{\text{vol}} \sim (\Delta_{\text{AD}}/H^{-1})^2$ for strings, if one assumes only one configuration per unit horizon. For fat walls and strings, a viable baryon number production is expected in the allowed parameter regions (10).}

We must also include the temporal A-term,

$$V_A(\phi, \phi^\dagger) = \left[ \frac{a'}{4M_p} \phi^4 + h.c. \right]$$ (1.7)

which disappears after $t > t_A$. Here $t_A$ denotes the time when the phase transition of the A-term takes place, and $a'$ is induced by the false vacuum configuration during $t < t_A$. $a'$ is required to be larger than $a_o$. Integrating the equation one can find the produced baryon number density in the domain of the false vacuum,

$$n_B \sim \frac{4q_B}{9H} \left| \frac{a_o}{M_p} \phi_{\text{AD}}^{\text{ini}} \right|^4 \delta_{\text{eff}}.$$ (1.8)

Unlike the conventional Affleck-Dine baryogenesis, here the Hubble parameter $H$ is the parameter that should be determined by the time of the A-term phase transition at $t = t_A$.

In Section 2 we show example in which the condensate is realized inside the domain of the quasi-degenerated vacua.

\footnotetext[3]{See eq.(1.5) in the review part.}
In section 3 we consider examples in which the misalignment of the A-term appears at the phase transition of the A-term.

2 Condensate in the false vacuum domain

In general models of Affleck-Dine baryogenesis, the destabilization of the flat direction is expected to be induced by the Hubble parameter, which is expected to be larger than \( m_{3/2} \) at an earlier stage of the Universe.

In this section we consider examples in which the Affleck-Dine condensate is realized inside the domain of the false vacuum. The most significant difference from the conventional scenario is that the condensate is realized in the local domain and it can remain at much lower temperature after reheating.

We consider an additional sector that is described by the superpotential

\[
W_M = \frac{1}{3} \lambda S^3 - \kappa_s S \Lambda^2 + \kappa_Q \bar{Q}^i Q^i. \tag{2.1}
\]

The quark-like superfield \( Q^i \) is not a conventional ingredient of the standard model, but charged with the standard model gauge group and \( U(1)_B \). \( Q^i \) can be replaced by the lepton-like superfield \( L^i \) if one considers Affleck-Dine leptogenesis. The minima of this potential are at

\[
\langle S \rangle = \pm \sqrt{\frac{\kappa_s}{\lambda}} \Lambda \\
\langle \bar{Q}^i Q^i \rangle = 0 \tag{2.2}
\]

and

\[
\langle S \rangle = 0 \\
\langle \bar{Q}^i Q^i \rangle \sim \Lambda^2. \tag{2.3}
\]

We set the scale \( \Lambda \) at the intermediate scale \( M_{EW} \ll \Lambda \ll M_p \), and assume that \( \kappa_s, \kappa_Q \) and \( \lambda \) are constants of \( O(1) \).

In this case, the Affleck-Dine field is not intended to be the flat direction of the model. The Affleck-Dine field is the baryonic direction that develops non-zero vacuum expectation value in the false vacuum (2.3). Here \( N_f \geq N_c \) heavy quarks are required in
our messenger sector, which decouple in the true vacuum while they condensate in the false vacuum. We can define baryonic combinations $B^{ij_1...i_{Nc}} = Q^{i_1}...Q^{i_{Nc}}$ as in the usual models of supersymmetric QCD, which is schematically denoted by $\phi_{AD}$. Obviously, this direction is not the flat direction of the model. One can also add heavy leptonic fields $L^i$ in this sector.

Here we make a brief comment on the stability of the false vacuum and related domain walls. First, the degeneracy between (2.2) and (2.3) is obviously broken by the soft mass terms. In the presence of the soft mass terms, the difference between these vacua becomes $\epsilon \sim m^2_{soft}\Lambda^2$ if there is no miraculous cancellation, which is enough to explain the safe decay. We should also consider the degeneracy between $S = \pm \sqrt{\kappa \Lambda}$. In this case, one may gauge the relevant discrete symmetry, or expect destabilization by the effective terms induced by the gravitational interactions.

### 3 A-term phase transition

In this section we construct several examples of the A-term phase transition that induces the driving force in the angular direction of $\phi_{AD}$.

#### 3.1 Thermal phase transition

We shall first consider an example where the thermal effect triggers the phase transition at the critical temperature $T_c$.

The phase transition can be realized by imposing extra gauge symmetries such as $U(1)_{B-L}$ or other string-motivated $U(1)$ symmetries. These extended gauge symmetries are sometimes utilized to certificate the proton decay in many kinds of grand unification models. We can consider the case where the higher dimensional terms that break baryon (or lepton) number are suppressed by enhanced gauge symmetries, which is sometimes discussed in the light of string-derived GUT models. These symmetries are used to suppress naturally the unsafe $d = 4$ as well as the color-triplet mediated or gravity induced $d = 5$ proton-decay operators, which generically arise in unification models of

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Affleck-Dine baryogenesis with gauged $U(1)_{B-L}$ is discussed in ref. [11] from a different point of view.
supersymmetry. For example, in the models that contain $B - L$ as a symmetry, the dangerous $d = 4$ operators permitted by the standard model gauge symmetry are forbidden. However, such operators can in general appear through non-renormalizable operators if there exists a spontaneous breaking of the symmetry in the presence of non-zero vacuum expectation values of the fields that carry $B - L$. To eliminate the dangerous $d = 5$ operators, one may add other additional gauge symmetries or adjust the charge assignment of the $B - L$ breaking field. Of course one may use discrete gauge symmetries alternatively\[14\].

Here we consider an additional superfield $Z$ charged with $U(1)_{B-L}$. We assume that $Z$ is a standard model singlet with the flat potential destabilized by the soft mass. We consider a non-renormalizable operator that contains $Z$,

$$W_A \sim \frac{1}{M_p^{n-3}} \left( \frac{Z}{M_p} \right) \phi^n_{AD}$$

(3.1)

where $\phi_{AD}$ denotes the AD direction that carries non-zero $U(1)_{B-L}$ charge. We have set that the cut-off scale is simply given by the Planck scale. The relevant A-term takes the following form

$$V_A \sim a m^{3/2} \frac{1}{M_p^{n-3}} \left( \frac{Z}{M_p} \right) \phi^n_{AD} + h.c. \quad (3.2)$$

where $a$ is the constant of $O(1)$.

$Z$ can be thermalized to obtain the thermal mass $\sim T^2$, which stabilizes the potential. Let us assume that the above superpotential dominates the A-term in the true vacuum. On the other hand, in the false vacuum where $Z$ vanishes, other terms can naturally contribute to determine the phase of the condensate so that the expected value of the phase becomes different from the one in the true vacuum.\[\square\] In this respect, the gap of the phase appears between the above two vacua. When the false vacuum of the A-term eventually collapses, the misalignment of the phase of $\phi_{AD}$ produces net baryon number in the local domain of the condensate $\phi_{AD} \neq 0$. In this model, the phase transition from $< Z > = 0$ to $< Z > \simeq \Lambda$ takes place at the temperature about $T_c \sim m^{3/2}$ since

\[7\] Without loss of generality, we can assume that there are other sources of A-term that determine the phase of the condensate when $Z$ vanishes. For example, there is the A-term of the form $V_A \sim \frac{F_Z}{M_p^3} \phi^n_{AD} + h.c.$ where $F_Z \sim \rho_I^{1/2}$ during thermal inflation.
the direction is expected to be flat around the origin. In this case, there should be a short period of thermal inflation, and the baryon number production occurs after thermal inflation. We can estimate the baryon to entropy ratio

$$\frac{n_B}{s} \sim \frac{T_{R_Z} |am_{3/2} Z \phi_{AD}^n|}{H_o M_p^{n-2} \rho_I}$$

(3.3)

where \( \frac{n_B}{n_{\phi AD}} \) is the ratio of the baryon to Affleck-Dine field number densities, which can be assumed to be \( \sim O(1) \). \( H_o \) denotes the Hubble parameter at the time when the oscillation of the AD field starts, and \( T_{R_Z} \) is the reheating temperature after thermal inflation. Assuming that the energy density during thermal inflation \( \rho_I \) is \( \rho_I \sim \left( \frac{m_{3/2}^3}{2 Z^2} \right) \), we can obtain the crude estimation for \( n = 4 \),

$$\frac{n_B}{s} \sim 10^{-10} \left( \frac{T_{R_Z}}{1 \text{GeV}} \right) \left( \frac{10^2 \text{GeV}}{m_{3/2}} \right)^2 \left( \frac{10^{12} \text{GeV}}{< Z >} \right)^2 \left( \frac{\phi_{AD}}{10^0 \text{GeV}} \right)^4,$$

(3.4)

where \( < Z > \) is the vacuum expectation value of \( Z \) in the true vacuum. Although in general \( H_o \) is expected to be less than \( H_I = \frac{\rho_Z^{1/2}}{M_p} \), here we have assumed \( H_o \approx H_I \) which corresponds to the case where the AD oscillation starts immediately after thermal inflation.

In this case, the baryon number is actually generated.

Let us discuss the string-derived symmetries for another example. Without any additional symmetry other than the standard model gauge symmetry, a supersymmetric theory in general permits dimension 4 and dimension 5 operators that violate the baryon and lepton numbers. The operators in problem are written by using standard notations,

$$W = [\eta_1 U D D + \eta_2 Q L D + \eta_3 L L E] = \frac{1}{M} [\lambda_1 QQQL + \lambda_2 U U D E + \lambda_3 LLH_2H_2].$$

(3.5)

Experimental limits on proton lifetime impose the constraint on the couplings as: \( \eta_1 \eta_2 \leq 10^{-24} \) and \( \lambda_{1,2}/M \leq 10^{-25} (\text{GeV})^{-1} \). One may also impose the requirements that the symmetry naturally ensures the proton stability and simultaneously permits the light neutrino masses. The solution given in ref.\( \text{[12, 13]} \) utilizes the gauged \( B - L \) and other string-motivated symmetries, such as

$$\lambda_1 \sim \left( \frac{\tilde{N}_R}{M} \right) \left( \frac{T_i T_j}{M^2} \right)^2 \sim 10^{-2.5} \left( \frac{\Lambda_c}{M} \right)^4,$$

(3.6)
where $T_iT_j$ is the condensate in the hidden sector and their magnitude is denoted by $\Lambda^2_c$. The experimental bound on $\eta_{1,2}$ implies $(\Lambda_c/M)^4 \leq 10^{-9.5}$. On the other hand, $\Lambda_c$ cannot be less than $T_{eV}$, since then the heavy particles of the relevant gauge symmetry would be accessible to the experiments.

Here we consider another operator

$$W_A = \lambda_A \frac{T_iT_j \phi^4_{AD}}{M_p M^2_p}.$$ 

(3.7)

In this model, naive thermal phase transition in the hidden sector ($T_iT_j = 0$ to $T_iT_j = \Lambda^2_c$) is expected to induce the misalignment of the phase. In this case, the temperature of the baryon number production is determined by the scale of the condensate of the hidden sector matter field $T_i$. The resultant baryon number is

$$n_B \sim \frac{1}{H} \left[ m_{3/2} |T_iT_j \lambda_A \phi^4_{AD}| \right]$$

(3.8)

where we have omitted numerical factors and parameters of $O(1)$. We estimate the baryon to entropy ratio at the critical temperature $T_c \sim \Lambda_c$,

$$\frac{n_B}{s} \sim 10^{-10} \left( \frac{\phi_{AD}}{10^{10} GeV} \right)^4 \left( \frac{10^6 GeV}{\Lambda_c} \right)^3 \left( \frac{m_{3/2}}{10^2 GeV} \right).$$

(3.9)

### 3.2 Less trivial example

Here we consider different types of phase transitions. In both cases, misalignment of the phase is induced by the changes in the effective A-term. The difference is that in the above examples the changes are induced by thermal phase transition while in the following examples they are consequences of the decay of the false vacuum domains of the A-term. In general, there are two types of such phase transitions. In one case, an alternative A-term becomes larger than the ordinary one, while in the other case the misalignment appears because the phase of the conventional A-term takes different values in the quasi-degenerated domains.

Larger A-term appears in the false vacuum

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9 See ref.[12] for notations.
10 See eq.(1.8) in the introduction.
The authors of ref.[17] argued that additional A-terms possibly appear by thermal effect and it will be proportional to the temperature $T$. In ref.[7], however, it is discussed that the A-term will not appear in generic situation.

Here we consider an alternative of A-term generation in ref.[17]. We consider the case where the condensate of the trigger field is not due to the thermal effect, but is realized by the false vacuum field configuration. We consider terms in the potential of the form:

$$ W = \lambda \frac{\phi^4_{AD}}{M} + h \phi_{AD} \xi \xi, $$  \hspace{1cm} (3.10)

where the field $\xi$ represents extra matter field that couples to the Affleck-Dine field with a Yukawa coupling. Higher dimensional couplings may be included in the coupling constant $\lambda$ and $h$. Then the potential includes terms such as:

$$ V_A = h \lambda \frac{\phi^3_{AD}}{M} \xi^* \xi^*. $$  \hspace{1cm} (3.11)

If large vacuum expectation value of $< \xi^* \xi^* >$ is formed in the false vacuum, quite large A-term will effectively appear. As we have discussed in the previous section, it is possible to construct such configurations in which the scalar component of the field $\xi$ develops large expectation value in the false vacuum. When the false vacuum bubble shrinks or collapses due to the energy difference, misalignment will appear.\footnote{Here the false vacuum of the A-term does not cover the whole Universe, but it appears as a false vacuum domain of the $H^{-1}$ size. Baryogenesis is expected to take place in the regions where both the field condensate and the additional A-term exist.} We can estimate the produced baryon number,

$$ n_B \sim \frac{1}{H} \left[ \lambda \frac{m^{3/2}}{M_p} \frac{\phi^4_{AD}}{M} \right], \hspace{1cm} (3.12) $$

where numerical factors and parameters of $O(1)$ are ignored. The resultant baryon to entropy ratio is

$$ \frac{n_B}{s} \sim 10^{-10} \left( \frac{\lambda}{10^{-11}} \right) \left( \frac{m^{3/2}}{10^2 GeV} \right) \left( \frac{\phi_{AD}}{10^6 GeV} \right)^4 \left( \frac{10^5 GeV}{T_A} \right)^5 \hspace{1cm} (3.13) $$

where $T_A$ denotes the temperature when the condensate of $\xi$ decays.

For an alternative example, here we consider a model with the superpotential

$$ W = S(XX - \Lambda^2_X) + S^3. \hspace{1cm} (3.14) $$
Here we assume that \( S \) is charged with the discrete gauged R-symmetry. The two degenerated minima of this potential are

\[
\langle S \rangle \simeq \pm \Lambda_X
\]

\[
\langle XX \rangle = 0
\]

and

\[
\langle S \rangle = 0
\]

\[
\langle XX \rangle \simeq \Lambda_X^2.
\]

The nonrenormalizable operators in the superpotential can be written as

\[
W_A \simeq \frac{1}{M_p^{n-3}} \left( \frac{S}{M_p} \right)^n \phi_{AD}^n
\]

where \( \phi_{AD} \) denotes the AD direction. We have set that the cut-off scale is simply given by the Planck scale.

When the false vacuum of \( S = 0 \) eventually collapses, the misalignment of the AD phase appears producing net baryon number in the local domain of the AD condensate. In general the temperature of the baryogenesis is determined by the scale of \( \Lambda_X \), from the requirement that the cosmological domain walls must decay before they dominate the Universe.\(^\text{12}\) For \( n = 6 \), the produced baryon number is

\[
n_B \sim \frac{1}{H} \left[ \frac{m_{3/2}}{M_p} \frac{\Lambda_X^{6/2}}{M_p^{3/2}} \phi_{AD}^6 \right]
\]

where numerical factors and parameters of \( O(1) \) are ignored. The resultant baryon to entropy ratio is

\[
\frac{n_B}{s} \sim 10^{-10} \left( \frac{\Lambda_X}{10^9 GeV} \right) \left( \frac{m_{3/2}}{10^2 GeV} \right) \left( \phi_{AD} \right)^6 \left( \frac{10^4 GeV}{T_A} \right)^5
\]

where \( T_A \) denotes the temperature when the false vacuum domain of the A-term collapses.

**Misalignment of the A-term**

We can also consider another example in which the domain wall appears simply as the consequence of the spontaneously broken \( Z_N \) symmetry of the coefficient in the A-term, which may appear when the A-term is determined by the field with the effective

\(^{12}\) The constraint is discussed in appendix A.
$Z_N$ symmetry. Then the misalignment of the A-term appears as the consequence of the discrete phases of a field, for example $<X> \simeq \Lambda_X e^{\frac{2\pi i k}{N}}$. The appearance of such misalignment is very natural since each field appearing in the A-term may have their effective discrete symmetries that is spontaneously broken by their vacuum expectation values.

In this case, there is a naive concern that the produced baryon in each $k$-domain may cancel when the contributions are summed over the whole Universe. To see what happens in our model, let us consider the relevant part of the potential of the form

$$V_A \sim \frac{\lambda a_k}{M_p^{n-3}} \phi^n_{AD} + h.c.$$  \hspace{1cm} (3.20)

where $a_k$ is assumed to be determined by the vacuum expectation value of the field. Here we consider the case where the misalignment occurs because the discrete domains of $a_k = \Lambda_a e^{\frac{2\pi i k}{N}}$ are formed before baryon number production. In addition to the domains of $k$-vacuum of the A-term, there are discrete $n$ vacua that appear because of the effective $Z_n$ symmetry of the AD field. As the result, the initial phase in each domain takes the form

$$\theta^{\text{ini}}_{\phi_{AD}} \simeq \frac{1}{n} \left[-\arg(a) - \arg(a_k) + 2\pi j\right],$$  \hspace{1cm} (3.21)

where $j = 1, 2, ..., n$ and $k = 1, 2, ..., N$. In the conventional Affleck-Dine baryogenesis, one of the $(j, k)$ vacuum dominates the whole Universe because of the inflationary period, which determines the initial condition of Affleck-Dine baryogenesis. On the other hand, in our case, the resultant baryon number may contain the following summation factor. Let us denote the A-term in the final state by $a_{k_0}$. If there is no biasing, it is given by

$$\epsilon_{\text{sum}} = \sum_{\text{All Domain}} \sin \left(\arg(a) + \arg(a_{k_0}) + n \arg(\phi_{AD})\right)$$

$$= \sum_{k=1}^{N} \sin \left(\arg(a) + \arg(a_{k_0}) + [-\arg(a) - \arg(a_k) + 2\pi j]\right)$$

$$= \sum_{k=1}^{N} \sin \left(\arg(a_{k_0}) - \arg(a_k)\right)$$

$$= 0$$  \hspace{1cm} (3.22)

In the above example, cancellation occurs if there is no biasing at the domain formation. On the other hand, however, complete cancellation is rather miraculous because

\footnote{The mechanism for biasing the domain formation is discussed by many authors. See ref.\cite{18} for}
the biasing is expected in more realistic situations. We expect the biasing factor of 
\( \epsilon_{\text{sum}} \sim 0.1 - 0.01 \) because the vacua as well as the potential is not strictly symmetric in our model.

Note that the cancellation cannot occur when the initial AD phase is determined by other A-terms.

4 Conclusions and Discussions

In this paper we have tried to separate the mechanism of Affleck-Dine baryogenesis from the inflation in order to avoid several obstacles in the original model. We have shown several examples for realizing Affleck-Dine baryogenesis in the alternative way. Our mechanism is quite efficient in producing baryons and we believe that application of our mechanism may solve the problems of the conventional models of baryogenesis. The idea of Affleck-Dine baryogenesis proceeding with different efficiency in local domains was suggested in ref. [19] from different context.

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A Constraint on the cosmological domain walls

It is well known that whenever the Universe undergoes a phase transition associated with the spontaneous symmetry breaking, domain walls will inevitably form. In most cases, the domain walls are dangerous for the standard evolution of the universe [8]. However, if some criteria are satisfied, unstable domain walls that disappear at \( t < t_c \sim (G\sigma)^{-1} \) can exist. First, we review how to estimate the constraint to safely remove the cosmological domain walls [9]. When the discrete symmetry is broken by gravitational interactions, estimation of the biasing factor requires computational simulations, which is not suitable for our motivation for this paper. In models for baryogenesis, the biasing factor is sometimes related to the CP violating phase of \( O(0.01) \).
the symmetry is an approximate discrete symmetry. The degeneracy is broken and the 
energy difference \( \epsilon \neq 0 \) appears. Regions of the higher density vacuum tend to collapse, 
the corresponding force per unit area of the wall is \( \sim \epsilon \). The energy difference \( \epsilon \) becomes 
dynamically important when this force becomes comparable to the force of the tension 
\( f \sim \sigma/R_w \), where \( \sigma \) is the surface energy density of the wall and \( R_w \) denotes the typical 
side for the wall distance. For walls to disappear, this has to happen before they be-
come harmful. On the other hand, the domain wall network is not a static system. In 
general, the initial shape of the walls right after the phase transition is determined by the 
random variation of the scalar VEV. One expects the walls to be very irregular, random 
surfaces with a typical curvature radius, which is determined by the correlation length of 
the scalar field. To characterize the system of domain walls, one can use a simulation[20]. 
The system will be dominated by one large (infinite size) wall network and some finite 
closed walls (cells) when they form. The isolated closed walls smaller than the horizon 
will shrink and disappear soon after the phase transition. Since the walls smaller than the 
horizon size will efficiently disappear so that only walls at the horizon size will remain, 
their typical curvature scale will be the horizon size, \( R \sim t \sim M_p/g^{1/2}T^2 \). Since the energy 
density of the wall \( \rho_w \) is about 
\[ \rho_w \sim \frac{\sigma}{R}, \]  
(A.1) 
and the radiation energy density \( \rho_r \) is 
\[ \rho_r \sim g_s T^4, \]  
(A.2) 
one sees that the wall dominates the evolution below a temperature \( T_w \) 
\[ T_w \sim \left( \frac{\sigma}{g_s^{1/2}M_p} \right)^{1/2}. \]  
(A.3) 
To prevent the wall domination, one requires the pressure to have become dominant before 
this epoch, 
\[ \epsilon > \frac{\sigma}{R_{wd}} \sim \frac{\sigma^2}{M_p^2}, \]  
(A.4) 
which is consistent with the criterion in ref.[3, 4]. Here \( R_{wd} \) denotes the horizon size at the 
wall domination. A pressure of this magnitude would be produced by higher dimensional 
operators, which explicitly break the effective discrete symmetry[16, 21].
The criterion (A.4) seems appropriate, if the scale of the wall is higher than \((10^5 GeV)^3\). For the walls below this scale \((\sigma \leq (10^5 GeV)^3)\), there should be further constraints coming from primordial nucleosynthesis. Since the time associated with the collapsing temperature \(T_w\) is \(t_w \sim M_p^2 / g^2 \sigma \sim 10^8 \left(\frac{(10^2 GeV)^3}{\sigma}\right)\) sec, the walls \(\sigma \leq (10^5 GeV)^3\) will decay after nucleosynthesis\(^{[22]}\). In this case, one must consider stronger constraint. If the walls are not hidden and are supposed to decay into standard model particles, the entropy produced when walls collapse will violate the phenomenological bounds for nucleosynthesis. On the other hand, this simple bound from the nucleosynthesis is not effective for the walls that cannot decay into standard model particles. The walls such as soft domain walls\(^{[23]}\), the succeeding story should strongly depend on the details of the hidden components and their interactions. These walls can decay late to contribute to the large scale structure formation.

Of course, the condition for the cosmological domain wall not to dominate the Universe (A.4) should also be changed if the wall velocity is lower than the speed of the light and then the Universe contains walls more than one. This implies that the condition to evade the wall domination becomes \(\epsilon > (\sigma^2 / M_p^2) \times x\), where the constant \(x\) is determined by \(R_w\) as \(x \approx M_p / (R_w T^2)\). For the walls with lower velocity, the bound for \(\epsilon\) is inevitably raised since such walls will dominate earlier.

B Comment on the early oscillation

In conventional scenarios of Affleck-Dine baryogenesis, the oscillation of condensates along flat directions that carry non-zero baryonic charge and attained large vacuum expectation values at the end of the inflationary epoch is important. Although these directions are flat, they may couple to other fields that acquire masses induced by the vacuum expectation values of the flat directions. If the masses of other fields are sufficiently small, the plasma of inflaton and their decay products can act on the flat directions, which may induce thermal masses. In such cases, the flat directions acquire large masses and start their oscillations earlier than usually estimated. There will be other way terminating the oscillations due to evaporation of the condensate. These thermal effects should alter the estimates of the conventional Affleck-Dine baryogenesis as is discussed in ref.\(^{[1, 17]}\).
In our paper we have discussed the phase transition of the A-term at low energy scales. However, the scales of the phase transitions can be raised to much higher energy scales. If the typical scales of such A-term phase transitions are raised so that they can take place before the early oscillation, they can produce baryon number before the problematic early oscillation avoiding the problem of the original Affleck-Dine baryogenesis\[10].

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