Electroweak baryogenesis mediated by locally supersymmetry-breaking defects

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We consider the scenario of electroweak baryogenesis mediated by a cosmological defect in models of supersymmetry breaking. When the effective electroweak breaking scale is raised in the defect configuration, the mechanism of electroweak baryogenesis works at a higher energy scale. The baryon charge produced by the mechanism is captured in the defect. It is protected from sphalerons and then released when the defect decays.

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I. INTRODUCTION

Contrary to a naive cosmological expectation, evidence shows that the Universe contains an abundance of matter over antimatter. In this paper we consider alternative mechanisms of electroweak baryogenesis. Electroweak baryogenesis is commonly known as an attractive idea because of its calculability in which testable physics, present in the standard model of electroweak interactions and its modest extensions, is responsible for this fundamental cosmological event. One may think that the previous negative results on the minimum standard model is an indication that the baryon number asymmetry in the Universe was not created at the electroweak phase transition, but rather is related to the physics of higher energy scales. Of course one can stick to electroweak baryogenesis considering the extensions of the particle content of the standard model to get a stronger electroweak phase transition in the allowed parameter region. In the general scenario of electroweak baryogenesis the baryogenesis occurs at the phase boundary; this requires the coexistence of regions of large and small $\langle H/T\rangle$, where H denotes the Higgs field in the standard model. In the regions of small $\langle H/T\rangle$, sphalerons are unsuppressed and can mediate baryon number violation, while the regions of large $\langle H/T\rangle$ are needed to store the created baryon number. Below the critical temperature T_c^{EW} of the electroweak phase transition, $\langle H/T\rangle$ grows till sphalerons are shut-off in the whole Universe. For electroweak baryogenesis to be possible, one needs some specific regions where $\langle H\rangle$ is displaced from the equilibrium value.

The idea we examine in this paper is that the same mechanism of electroweak baryogenesis can happen along topological defects left over from some other cosmological phase transitions that took place before the electroweak phase transition. The idea of defect-mediated electroweak baryogenesis is already discussed by many authors [1]. They have considered the configurations of the Higgs field itself where the vanishing Higgs vacuum expectation value is realized in the core of the cosmological defects. On the other hand, the displaced Higgs vacuum expectation value can be obtained by indirect effects of other field configurations. If the effective electroweak symmetry-breaking scale is raised in some regions inside the cosmological defect, sphalerons could be suppressed in such regions while they would be effective in the bulk of space. The motion of the defect network, in a similar way as the motion of bubble walls in the usual strongly first-order phase transition scenario, will leave a net baryon number behind the moving surface and then the baryon asymmetry will be kept in the sphaleron-suppressed regions inside the defect. Then the defect protects the baryon charge from sphalerons until sphalerons become inactive, and then decays to release the baryon number after electroweak phase transition. The idea of such protection of the produced baryon number is not new and is commonly used by many authors in the topics such as Q balls [2] or defect-mediated baryogenesis [3,4].

In this paper we will point out that this idea works in supersymmetric extensions of the standard model when the supersymmetry-breaking scale is raised inside the defects. The defects should be formed before the electroweak phase transition and should decay after the electroweak phase transition. We consider the defects that break supersymmetry locally and only at a short period of cosmological evolution of the Universe.

In Ref. [4] the mechanism of baryogenesis by the decay of the string defects which initially possess the baryon number is discussed. The mechanism we will discuss in our paper can be one of the possible mechanisms of generating the initial baryon number asymmetry contained in the defect.

II. GAUGE-MEDIATED SUPERSYMMETRY BREAKING IN THE DEFECTS

Many kinds of mechanisms of breaking supersymmetry are discussed by many authors. The hidden sector model in supergravity [5] is perhaps one of the most popular scenarios among them. In hidden sector models, supersymmetry is broken in the hidden sector by some mechanisms, such as the Polonyi model [6], dynamical breaking by gaugino condensation [7], or the O'Raifeartaigh model [8]. The effects of the supersymmetry breaking are mediated to the fields in the supersymmetric standard model only by the gravitational interactions, thus suppressed by the cut-off scale.

There is another mechanism of supersymmetry breaking in which the effects are mediated by gauge interactions. The gauge mediation of supersymmetry breaking has an attractive feature, ensuring the degeneracy of squark masses and therefore suppresses the dangerous flavor changing neutral current (FCNC) effects.

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The main motivation for these models is to explain the origin and the stability of the hierarchy between the fundamental scale and the electroweak scale. In this sense, the quadratic divergence in the Higgs boson mass parameter coming from a top quark radiative correction is canceled by the divergence coming from a scalar top. Considering these cancellations and including supersymmetry breaking at scale m_{SUSY}, the resulting divergence becomes logarithmic. The Higgs mass is reliably computed in the effective theory, and is not dominated by unknown physics at the cutoff. To be more precise, one can say that the stability of the hierarchy is due to the existence of supersymmetry at a higher energy scale, while the hierarchy is produced by the dynamical breaking of the supersymmetry or large suppression factor from the fundamental scale. One can also find some alternatives of these mechanisms of supersymmetry breaking in light of the M theory, large extra dimensions, and brane worlds [9].

Among these mechanisms of supersymmetry breaking, we first focus our attention on the gauge-mediated models. In this section we explore the possibility of obtaining the baryon asymmetry of the Universe by using toy models of gauge-mediated supersymmetry breaking.

A. Toy model 1

Here we would like to consider the simplest version of the gauge-mediated models of supersymmetry breaking which does *not* break supersymmetry in the true vacuum but breaks supersymmetry locally in the defect configuration. The messenger sector can be described by the superpotential

$$W_M=\frac{1}{3}\lambda S^3-\kappa_s S\Lambda^2+\kappa_q S q\bar{q}+\kappa_l S l\bar{l}, \tag{2.1}$$

where S is a singlet superfield. The superfield q transforms as a $(3,1,\frac{1}{3})$ under the standard model, while l transforms as $(1,2,-\frac{1}{2})$. Then the minima of this potential are at

$$\begin{aligned}\langle S\rangle&=\pm\sqrt{\frac{\kappa_s}{\lambda}}\Lambda,\\ \langle\bar{q}q\rangle&=\langle\bar{l}l\rangle=0\end{aligned} \tag{2.2}$$

and

$$\begin{aligned}\langle S\rangle&=0,\\ \langle\bar{q}q\rangle&\sim\langle\bar{l}l\rangle\sim\Lambda^2.\end{aligned} \tag{2.3}$$

We set the scale Λ at the intermediate scale $\Lambda\sim10^6$ GeV. κ_s, κ_q, and κ_l are assumed to be $O(1)$. λ is not assumed to be $O(1)$, since the width of the domain wall is determined by λ. Two examples are considered, $\lambda\sim10^{-1}$ (thin wall) and $\lambda\sim10^{-6}$ (fat wall).

At the boundary of $\langle\bar{q}q\rangle=\langle\bar{l}l\rangle=0$, it is easy to see that the auxiliary component of the field S (F_S) can acquire vacuum expectation value of order $\langle F_S\rangle\simeq\Lambda^2$ in the wall. This effect feeds down to the maximal supersymmetric standard model (MSSM) sector through loop corrections. The soft terms calculated this way depend on the parameter $\langle F_S/S\rangle$ in the messenger sector. In this model we assume that the effective electroweak scale in this region becomes as large as 10^4 GeV.

If the second vacuum (2.3) is lifted by a small soft mass of q and l, or by contributions suppressed by the cut-off scale, the degeneracy is broken and the domain of the local minimum (2.3) shrinks. It occurs when the pressure ϵ dominates the energy density and also becomes larger than the force of the wall tension.[1] Denoting the induced soft mass as m_0, the temperature when the domain shrinks is $T_s\simeq\sqrt{m_0\Lambda}$. For $m_0>10^2$ GeV, the domain shrinks before baryogenesis starts, and the wall appears as the composite state during baryogenesis. The composite domain wall is constituted by two parts, the outer region $|S|>|F_S|^{1/2}$, where the messenger matter fields q and l acquire large mass and stay at the origin, and the inner region $|S|<\Lambda$, where the messenger matter fields q and l can develop nonzero vacuum expectation values. The outer side of the wall behaves almost the same as the conventional gauge mediation sector of supersymmetry breaking and works as the phase boundary in the conventional defect-mediated electroweak baryogenesis, producing the required gap. In the inner region, which we denote the "core" of the defect, the messenger matter fields q and l can develop nonzero vacuum expectation values and protect the incoming baryon asymmetry produced at the outer side of the wall.

When m_0 is smaller than the electroweak scale ($m_0<10^2$ GeV), the false vacuum $\langle S\rangle=0$ appears during the baryogenesis, and then it shrinks at $T=T_s$. This period corresponds to the limit when the width of the core becomes infinite. In the second vacuum configuration (2.3), where the gauge symmetry is already broken by nonzero expectation values of q and l, the induced soft mass m_0 must be small if the supersymmetry breaking in the bulk of space is induced by an another sector of gauge-mediated supersymmetry breaking.

Here we should note that the Higgs field is not the sole candidate of the sphaleron-suppressing field that condensates inside the baryon-protecting defect. It is easy to see that any field that carries the $SU(2)_L$ quantum numbers contributes to the sphaleron energy, so that the baryon number breaking sphaleron interactions can be suppressed by the field condensates of other fields in the core [4]. We should also note that if a condensate is carrying the baryon number, then there will be massless excitations of the Goldstone boson as well. What we consider is the situation that the baryon number is spontaneously broken and the baryonic charge is stored inside the defect. The situation is very similar to the well-known idea of B balls or the baryogenesis discuused in Ref. [4].

What we will consider is the situation when the effective scale of the soft supersymmetry-breaking parameter $\langle|F_S/S|\rangle$ is raised in the surface region of the defect, but the particle spectrum is not affected in the bulk of space. This

[1]Here ϵ denotes the energy difference between two minima, Eqs. (2.2) and (2.3).

can be realized in a simple way if the dynamical supersymmetry-breaking sector or the messenger sector develops a cosmological defect. In the simplest case (2.1) the excessive breaking of supersymmetry is realized in the (composite) cosmological domain wall which interpolates two minima $\langle S\rangle = \pm\sqrt{(\kappa_s/\lambda)}\Lambda$. The supersymmetry breaking in the defect sector can vanish in the true vacuum, but should become large in the defect to realize the coexistence of the regions of large and small $\langle H/T\rangle$.

Since the defect sector is not necessarily required to be responsible for the soft terms in the MSSM in the true vacuum, there are no complexities related to the dynamical breaking of supersymmetry at the global minimum, the constraint on the CP breaking parameter, etc.[2] In general, the electroweak scale is intimately related to the soft-breaking parameters which can be raised in the ''local'' region in the defect.[3] At the temperature $T_c > T > T_{EW}$, baryon asymmetry produced in front of the defect can be trapped in the defect. Defects are able to trap the baryon from the time of the electroweak symmetry-breaking phase transition in the defect ($T = T_c \simeq 10^4$ GeV) till the Universe cools down to $T = T_{EW}$. Then the defects release the baryon number and finally disappear at $T = T_d$.

Here the mechanism of baryon asymmetry generation itself is similar to the conventional defect-mediated electroweak baryogenesis. Historically, the ways in which baryons may be produced when a phase boundary sweeps through space have been separated into two categories. One is called ''local baryogenesis'' in which baryons are produced when the baryon number violating processes (sphaleron interactions) and the CP violating processes induced by the wall occur together near the bubble walls, and the other is called ''nonlocal baryogenesis'' in which particles undergo CP violating interactions with the bubble wall and then become the flux of an asymmetry which carries a quantum number other than the baryon number into the unbroken phase region away from the wall. In the latter case baryons are then produced as baryon number violating processes convert them into the asymmetry in the baryon number. In general, both of these two ways of baryogenesis will occur together and the baryon number asymmetry of the Universe will be expressed by the sum of that generated by the two coexisting processes. When the speed of the phase boundary is greater than the sound speed in the plasma, local baryogenesis will dominate. Otherwise, nonlocal baryogenesis is usually more efficient.

Let us consider the simplest case of conventional nonlocal baryogenesis, and then examine the electroweak baryogenesis induced by the simplest supersymmetry-breaking defect. When the thin boundary limit is considered, the final baryon to entropy ratio of the conventional electroweak baryogenesis becomes [3]

$$\frac{n_B}{s} \sim 0.2\alpha_W^2 (g^*)^{-1} \kappa \Delta\theta_{CP} \frac{1}{v_w}\left(\frac{m_l}{T}\right)^2 \frac{m_h}{T}\frac{\xi^L}{D_L}, \qquad (2.4)$$

where D_L is the diffusion constant for leptons, and ξ^L is the persistence length of the current in front of the bubble wall. Here m_l and m_h denote the lepton and Higgs boson masses. When the background field configuration is steep, at the temperature much below the phase transition in the defect at $T = T_c$, the effect of CP violation is suppressed exponentially since the typical energy of the charge carrier is lower than the potential barrier. In this sense the formula (2.4) can be applied in the case that the energy of the leptons is comparable to the Higgs vaccum expectation value (VEV) inside the phase boundary. In our simplest case (2.1), the width of the defect depends on the parameter λ. When $\lambda \simeq O(0.1)$ the width of the domain wall is simply given by $\Delta \sim \Lambda^{-1}$, and the background defect configuration is steep. In a conventional scenario of electroweak baryogenesis, the flux is injected by the phase boundary into the unbroken phase and it is converted into the baryon asymmetry in the unbroken phase near the phase boundary. Then the produced baryons are trapped in the broken phase. In this respect, the mechanism of baryogenesis in our model is similar to the conventional mechanism of electroweak baryogenesis.

The wall which interpolates between Eqs. (2.2) and (2.3) can be divided into two parts. In the outer half of the wall, the supersymmetry-breaking parameter is locally raised so that the effective scale of the electroweak symmetry-breaking increases simultaneously. Considering the conventional calculation of the injected flux [10], one can confirm that the injected flux from the outer half of the domain wall into the sphaleron-activated region induces the baryon asymmetry which is very similar to the conventional electroweak baryogenesis (2.4). One may worry that the injected flux from the inner half may cancel the one from the outer half. Of course, it is hard to believe that they cancel exactly even if the alternating signs of the same magnitude of injected fluxes are expected by the naive order estimation. Moreover, the contribution to the injected flux from the inner half is strongly model dependent. It depends not only on the defect sector, but also on what one chooses for the mechanism of supersymmetry breaking. The important point is (1) when the magnitude of the injected flux from the outer half is larger than the one from the inner half, Eq. (2.4) with larger mass scales dominates and (2) when the magnitude of the injected flux from the inner half is larger than the one from the outer half, it induces larger baryon asymmetry. In this case, one can arrange the parameters or add new fields in the defect sector to obtain the desirable result. These models will be interesting, but should be discussed in another place since these issues do not match our motivations of this paper. Here we do not discuss the case (2) any longer.

On the other hand, when λ is quite small ($\sim 10^{-6}$), the background field configuration can be fat enough so that the effective soft supersymmetry-breaking mass is well approximated by a constant around the regions of electroweak baryogenesis, then the profile of the phase boundary is just the same as the one in the conventional electroweak phase transition. (Of course the effective-mass scales are higher

[2] We will comment on these issues in Appendixes.

[3] Of course, this naive expectation is not always correct. We will also comment on this relation in the Appendixes, focusing on the μ problem.

than the conventional one.) The condition of such a fat background is schematically given by the linear approximation as

$$m_h^{MAX}\times\left(\frac{T^{-1}}{\Delta}\right)\ll m_h^{PB}, \tag{2.5}$$

where m_h^{MAX}, Δ, and m_h^{PB} denote the maximum value of m_h inside the defect, the width of the wall, and the Higgs boson mass at the phase boundary. Let us consider the case where the defect has the thin phase boundary and the fat background, and that most of the baryon number is caught in the defect.[4] Then we should integrate n_B during the period of $T_c>T>T_{end}$, where T_{end} denotes the temperature when the fat background approximation breaks down and the suppression of the CP violation starts. One should also consider the effective volume that the defects sweep, whose suppression is negligible for the domain walls but can be important to strings.

When the background wall configuration is fat the baryogenesis lasts long, thus the baryon number asymmetry is enhanced. For fat walls, the energy gap that appears at the phase boundary is determined so that the velocity of the phase boundary equals with the speed of the background wall. It is easy to see that $\langle H\rangle/T$ as well as the energy gap increases as the phase boundary moves inside, while it decreases in the outer region. In this respect, the critical point where the phase boundary appears should depend on the velocity of the background domain wall. Here we postpone the analysis on peculiar situations ($v_w=1$ or $v_w\ll 1$) but consider the case when the velocity of the wall is close to the velocity of the conventional electroweak phase transition. Then m_l/T, m_h/T, and ξ^L/D_L in Eq. (2.4) are expected to be nearly the same as the conventional electroweak baryogenesis.

Taking these into account, we conclude that the baryogenesis mediated by the locally supersymmetry-breaking defects is a promising candidate for the baryon asymmetry of the Universe (BAU). The mechanism of baryogenesis itself is the same as that used in the conventional electroweak baryogenesis. The novel issue in this attempt is the origin of the defects that enables the electroweak baryogenesis at an earlier (and longer) period of the Universe when the wall is fat enough. The electroweak phase transition itself is not required to be first order, which is the same characteristic of the conventional defect-mediated electroweak baryogenesis.

B. Toy model 2

Let us consider another example of the gauge-mediated model of locally broken supersymmetry. Here we consider the dynamical supersymmetry breaking in the vectorlike gauge theories [11]. Denoting the singlet in the supersymmetry-breaking sector by Z, the low-energy effective superpotential is given by

$$W_{eff}=\lambda_Z\Lambda^2 Z, \tag{2.6}$$

where Λ is the dynamically generated scale in the supersymmetry-breaking sector. The effective Kähler potential is expected to take a form

$$K=|Z|^2-\frac{\eta}{4\Lambda^2}|\lambda_Z Z|^4+\cdots, \tag{2.7}$$

where η is a real constant of order one. Then the effective potential is given by

$$V_Z\simeq|\lambda_Z|^2\Lambda^4\left(1+\frac{\eta}{\Lambda^2}|\lambda_Z|^4|Z|^2\right). \tag{2.8}$$

If $\eta<0$ a nonzero vacuum expectation value of the singlet is expected. The F component of the singlet is nonvanishing, and it is expected to be $\langle F_Z\rangle\simeq\lambda_Z\Lambda^2$. The width of the defect configuration is $\Delta\simeq m_Z^{-1}\simeq(\eta^{1/2}\lambda_Z^3\Lambda)^{-1}\sim(10\text{ GeV})^{-1}$ for $F_Z^{1/2}=10^6$ GeV and $\lambda_Z=10^{-2}$. In this model, the fat defect can appear in a natural parameter region.

First we consider the case when a singlet Z couples directly to the messenger matter fields and a cosmological defect is formed for the singlet Z, which is charged under global $U(1)_R$. The spontaneous breakdown of the $U(1)_R$ symmetry in the earlier period of the Universe produces the global string network. The $U(1)_R$ symmetry, however, must be an approximate symmetry since it must be broken at least by an explicit breaking constant term in the superpotential in order to set the cosmological constant to zero. Such an explicit breaking term induces the domain wall configuration bounded by the string, which decays soon after it is formed. The energy difference that lifts the degeneracy in the $U(1)_R$ rotation is about $\epsilon_R\sim|W_{eff}|/M_p^2$ [12]. Then the $U(1)_R$ wall-string network decays when $T=\epsilon_R^{1/4}\sim 1$ GeV. The inner structure of the defect is almost the same as the domain wall in the toy model 1.

One can also consider another example where the defect is a local string and the configuration in the core breaks color symmetry, developing the squark vacuum expectation value. This assumption is natural, since the (unstable) color-breaking minimum is a natural feature of the supersymmetry breaking. Of course, one can introduce an additional defect sector that induces the required symmetry breaking, as we have discussed above. In this case, the baryon number is assumed to be broken spontaneously inside the string, and the baryonic charge may be stored in the core. Denoting the squarks as $\tilde{q}$, they carry a $U(1)$ baryonic global charge which is derived from the conserved current

$$J_B^\mu=\frac{i}{2}\sum_q q_B^q(\tilde{q}^\dagger\partial^\mu\tilde{q}-\tilde{q}\partial^\mu\tilde{q}^\dagger), \tag{2.9}$$

where q_B^q is the baryonic charge associated with any field $\tilde{q}$. Assuming the cylindrical symmetry, the baryonic charge per

[4]Although the background changes gradually in fat defects, the electroweak phase transition occurs at the critical point which moves as the temperature changes. The phase boundary can be much thinner than the background defect.

unit length (Q_B) along the z axis will be given by the integration $Q_B = \int d\theta dr\, r j_B(\theta,r)$. This type of string is expected to generate the suitable baryon number asymmetry of the Universe, if some conditions are satisfied. In light of Ref. [4], our mechanism of electroweak baryogenesis works to seed the initial baryon number confined in the string defects.

By interpolating two degenerated vacua in separate regions of space, one obtains a domain wall. If we have three or more discrete vacua in separate regions of space, segments of domain walls can meet at a one-dimensional "junction." These junctions can have a structure that is very similar to the strings. Although the evolution of the junctions is different from the strings and probably much more complicated to be analyzed, it seems possible to construct the model to produce the baryon asymmetry in the Universe.

One can also consider an alternative of the model in which the mass of the messenger matter field is produced by the expectation values of other fields instead of $\langle Z \rangle$ [13]. Then our mechanism works when the cosmological defect is formed by the fields that generate the mass terms for messenger matter fields q and l.

We also note that the baryons produced by other mechanisms before the electroweak phase transition can survive the wash out if they are trapped in the supersymmetry-breaking defects that we have discussed in this paper. This may also open another possibility for other baryogenesis.

III. CONCLUSIONS AND DISCUSSIONS

In this paper we examined new possibilities for electroweak baryogenesis mediated by cosmological defects. We analyzed the supersymmetric theories in which the hierarchy is produced by the soft breaking of supersymmetry. Although the magnitude of the baryon asymmetry depends on the profiles of the defects, the idea is general and can be applied to many models of supersymmetry. We also note that the baryons produced by other mechanisms before the electroweak phase transition can survive the wash out if they are trapped in the supersymmetry-breaking defects. This may open another possibility for other baryogenesis.

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APPENDIX A: CONSTRAINTS ON COSMOLOGICAL DOMAIN WALLS

It is well known that when the Universe undergoes a phase transition that is associated with the spontaneous symmetry breaking of discrete symmetries, domain walls will inevitably form. In most cases the domain walls must be removed since they are dangerous for the standard evolution of the universe. In this appendix we give a short review to show how to estimate the constraint to safely remove the dangerous cosmological walls. The crudest estimate we can make will be to insist that the walls are removed before they dominate over the radiation energy density in the Universe. When the explicit breaking of the discrete symmetry is expected because of the gravitational interactions, the symmetry must be an approximate symmetry. Then the degeneracy of the vacua is lost and the energy difference $\epsilon \neq 0$ appears. When ϵ dominates the energy density of the false vacuum, regions of the higher density false vacuum tend to shrink. The corresponding force per unit area of the wall is $\sim \epsilon$. The energy difference ϵ becomes dynamically important when this force becomes comparable to the force of the tension $f \sim \sigma/R_w$, where σ is the surface energy density of the wall and R_w denotes the typical scale of the wall distance. For walls to disappear safely, this has to happen before the walls dominate the Universe. 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 may expect the walls just after they are formed 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, simulations [14] are commonly used. According to the simulations, the system will be dominated by one large (infinite size) wall network and some finite closed walls (cells) just after the phase transition. The isolated closed walls smaller than the horizon will shrink and disappear soon after they are formed. 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, $R \sim t \sim M_p / g_*^{1/2} T^2$. Then the energy density of the wall ρ_w is about

$$\rho_w \sim \frac{\sigma}{R}, \tag{A1}$$

and the radiation energy density ρ_r is

$$\rho_r \sim g_* T^4, \tag{A2}$$

and one can see that the wall domination starts below a temperature T_w,

$$T_w \sim \left(\frac{\sigma}{g_*^{1/2} M_p} \right)^{1/2}. \tag{A3}$$

To prevent the wall domination, one requires the pressure ϵ to have become dominant before this epoch. This requires the constraint

$$\epsilon > \frac{\sigma}{R_{wd}} \sim \frac{\sigma^2}{M_p^2}. \tag{A4}$$

Here R_{wd} denotes the horizon size at the wall domination. A pressure of this magnitude is expected to be produced by higher-dimensional operators which explicitly break the discrete symmetry. The requirement is satisfied in general models of supersymmetry if the symmetry is a discrete R-symmetry Z_n^R [12].

In our model the requirement that the walls decay after electroweak symmetry breaking also imposes the upper bound on σ as $\sigma < (10^8\ \text{GeV})^3$, which excludes the hidden (M_p suppressed) sector for the defect.

Here we should note about the lower bound that comes from the nucleosynthesis. The criterion (A4) seems appropriate if the scale of the wall is higher than $(10^5\ \text{GeV})^3$. For the walls below this scale $[\sigma \leqslant (10^5\ \text{GeV})^3]$, there should be further constraints coming from primordial nucleosynthesis. Since the time associated with the collapsing temperature T_w is $t_w \sim M_p^2/g_*^{1/2}\sigma \sim 10^8\ [(10^2\ \text{GeV})^3/\sigma]$ sec, the walls $\sigma \leqslant (10^5\ \text{GeV})^3$ will decay after nucleosynthesis [15] and violate the phenomenological bounds for nucleosynthesis. If the walls are not hidden and can decay into the standard model particles, the entropy produced when walls collapse will violate the phenomenological bounds for nucleosynthesis. On the other hand, the succeeding story should strongly depend on the details of the hidden components and their interactions if the walls are soft domain walls [16]. They 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 (A4) should also be modified if the wall velocity is lower than the speed of light and if the Universe contains more than one wall. 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 $x \simeq M_p/(R_w T^2)$. For the walls with lower velocity, the bound for ϵ is inevitably raised since such walls will dominate earlier.

APPENDIX B: μ TERM AND CP VIOLATION

In order to make our discussions simple and generic, we made a naive assumption that the relative relations between mass parameters that appear in the conventional scenario for electroweak baryogenesis are not drastically changed at the surface of the defect so that one can use the conventional mechanism of electroweak baryogenesis. When the mechanism of generating the μ and B_μ terms in the defect is completely different from the one for the supersymmetry-breaking soft masses, the μ and B_μ terms are in general not altered and remain the same in the locally supersymmetry-breaking defects. Then the effective soft supersymmetry-breaking mass becomes 10–10^2 times larger than the μ and B_μ terms in the effective theory at the defect surfaces. One may also consider other alternatives in which the structure of the effective low-energy theory becomes completely different in the defect. These models will be interesting, but should be discussed an another time since these issues do not match our motivations for this paper.

Of course, one knows in some models the μ or B_μ terms can be related to the supersymmetry-breaking parameters F_X/X [17,18]. The most interesting case is that the μ term originates from the interaction with the supersymmetry-breaking sector, while the B_μ is suppressed so that the SUSY CP problem is solved in the bulk of space. In such models for μ and B_μ terms generation, almost all the input mass scales are determined by the supersymmetry-breaking parameter F_X/X and the mechanism of the electroweak baryogenesis at the defect surface looks precisely the same as the one for the conventional MSSM.

One can construct models in which the generating mechanism of B_μ in the defect is different from the one in the bulk of space, so that the effective theory at the defect surface can develop a large B_μ parameter. In this case, B_μ^{eff} locally induces a large CP parameter and the electroweak baryogenesis is enhanced.[5]

In our model for electroweak baryogenesis, one can expect another contribution to these CP violating phases from locally supersymmetry-breaking defects, since the origin of the supersymmetry-breaking mass and μ and B_μ parameters is completely or partially different from the true vacuum, changing the combination $\theta_{phys} \equiv \text{Arg}(\mu B_\mu^* m_G) = \theta_\mu - \theta_B + \theta_G$.

[5]In the gauge-mediated model, the phases in the gaugino masses m_G, μ, and B_μ parameters are physical, and in general, they can be large enough to conflict with experimental constraints. There are many ways to restrict these CP violating phases. For example, the SUSY CP problem can be solved if B_μ vanishes at the messenger scale, which is further discussed in Ref. [19].

[1] R. Brandenberger, A.C. Davis, and M. Trodden, Phys. Lett. B **335**, 123 (1994); Phys. Rev. D **53**, 4257 (1996); Phys. Lett. B **349**, 131 (1995); Phys. Rev. D **54**, 6059 (1996).

[2] A. Kusenko, Phys. Lett. B **404**, 285 (1997).

[3] A. Riotto and M. Trodden, Annu. Rev. Nucl. Part. Sci. **49**, 35 (1999).

[4] R. Brandenberger and A. Riotto, Phys. Lett. B **445**, 323 (1999).

[5] R. Barbieri, S. Ferrara, and C.A. Savoy, Phys. Lett. **119B**, 343 (1982); A.H. Chamseddine, R. Arnowitt, and P. Nath, Phys. Rev. Lett. **49**, 970 (1982); L.J. Hall, J. Lykken, and S. Weinberg, Phys. Rev. D **27**, 2359 (1983).

[6] J. Polonyi, Budapest Report No. KFKI-93 (1977).

[7] P. Niles, Phys. Lett. **115B**, 193 (1982); Nucl. Phys. **B217**, 366 (1983); U. Ellwanger, Phys. Lett. **153B**, 257 (1985); A. Font, L.E. Ibanez, D. Lust, and F. Quevedo, Phys. Lett. B **245**, 401 (1990); T.R. Taylor, *ibid.* **252**, 59 (1990); B. De Carlos and M. Moretti, *ibid.* **341**, 302 (1995); Z. Lalak, A. Niemeyer, and H.P. Nilles, Nucl. Phys. **B453**, 100 (1995); I. Antoniadis and M. Quiros, *ibid.* **B505**, 109 (1997).

[8] L. O'Raifeartaigh, Nucl. Phys. **B96**, 331 (1975).

[9] M.A. Luty and R. Sundrum, Phys. Rev. D **62**, 035008 (2000); Z. Chacko and A.E. Nelson, *ibid.* **62**, 085006 (2000); I. Antoniadis, C. Munoz, and M. Quiros, Nucl. Phys. **B397**, 515 (1993); A. Delgado, A. Pomarol, and M. Quiros, Phys. Rev. D **60**, 095008 (1999); D.E. Kaplan, G.D. Kribs, and M. Schmaltz, *ibid.* **62**, 035010 (2000); N. Arkani-Hamed, L. Hall, D. Smith, and N. Weiner, *ibid.* **63**, 056003 (2001).

[10] M. Joyce, T. Prokopec, and N. Turok, Phys. Rev. D **53**, 2930 (1996).

[11] K.-I. Izawa and T. Yanagida, Prog. Theor. Phys. **94**, 1105

(1995); **95**, 829 (1996); T. Hotta, K.-I. Izawa, and T. Yanagida, Phys. Rev. D **55**, 415 (1997).
[12] T. Matsuda, Phys. Lett. B **436**, 264 (1998); **486**, 300 (2000).
[13] K.-I. Izawa, Y. Nomura, K. Tobe, and T. Yanagida, Phys. Rev. D **56**, 2886 (1997).
[14] J.A. Harvey, E.W. Kolb, D.B. Reiss, and S. Wolfram, Nucl. Phys. **B201**, 16 (1982); T. Vachaspati and A. Vilenkin, Phys. Rev. D **30**, 2036 (1984).
[15] S.A. Abel, S. Sarker, and P.L. White, Nucl. Phys. **B454**, 663 (1995); T. Han, D. Marfatia, and Ren-Jie Zhang, Phys. Rev. D **61**, 013007 (2000).
[16] J.A. Frieman, C.T. Hill, and E. Watkins, Phys. Rev. D **46**, 1226 (1992); A. Singh, *ibid.* **50**, 671 (1994); G.M. Fuller and D.N. Schramm, *ibid.* **45**, 2595 (1992); A. Massarotti and J.M. Quashnock, *ibid.* **47**, 3177 (1993); S. Lola and G.G. Ross, Nucl. Phys. **B406**, 452 (1993).
[17] G. Dvali G.F. Giudice, and A. Pomarol, Nucl. Phys. **B478**, 31 (1996).
[18] T. Moroi, Phys. Lett. B **477**, 75 (1997); M. Dine, Y. Nir, and Y. Shirman, Phys. Rev. D **55**, 1501 (1997).
[19] R. Rattazzi and U. Sarid, Nucl. Phys. **B501**, 297 (1997); E. Gabrielli and U. Sarid, Phys. Rev. Lett. **79**, 4752 (1997).