Spontaneously broken parity and consistent cosmology with transitory domain walls

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(Received 13 November 2009; published 10 February 2010)

Domain wall structure which may form in theories with spontaneously broken parity is generically in conflict with standard cosmology. It has been argued that Planck scale suppressed effects can be sufficient for removing such domain walls. We study this possibility for three specific evolution scenarios for the domain walls, with evolution during a radiation dominated era, during a matter dominated era, and that accompanied by weak inflation. We determine the operators permitted by the supergravity formalism and find that the field content introduced to achieve desired spontaneous parity breaking makes possible Planck scale suppressed terms which can potentially remove the domain walls safely. However, the parity breaking scale, equivalently the Majorana mass scale M_R of the right-handed neutrino, does get constrained in some of the cases, notably for the matter dominated evolution case which would be generic to string theory inspired models, giving rise to moduli fields. One left-right symmetric model with only triplets and bidoublets is found to be more constrained than another admitting a gauge singlet.

DOI: 10.1103/PhysRevD.81.045010 PACS numbers: 12.10.−g, 11.27.+d, 12.60.Jv

I. INTRODUCTION

Existence of right-handed neutrino states [1–4] is a strong indicator of parity as a fundamental symmetry of nature, spontaneously broken in low energy physics. The scale of parity breakdown is as yet unknown. The seesaw mechanism [5–8] while providing an elegant qualitative explanation is unable to make a narrow prediction of the relevant energy scale due to wide variation in the fermion masses across the generations. In this paper we study the role of domain walls formed by spontaneous parity breaking in determining or constraining the energy scale of breakdown of this symmetry.

To be specific, we study the implementation of parity in left-right symmetric models [9–13]. While such models can descend from an $SO(10)$ grand unified theory [14–18], the scale of such unification is known to be high ($\sim 10^{16}$ GeV). On the other hand, the scale of parity breaking is completely undetermined and could be much lower. Indeed there are no observational obstructions for the parity breaking scale (and the associated right-handed neutrino Majorana mass scale)to be as low as TeV scale [19–21]. Here we study a particular source of constraint on this scale imposed by cosmology, with the possibility of restricting the parity breaking scale to low values. It suffices for this purpose to focus on the left-right symmetric model alone, independent of how the model may unify into $SO(10)$. For the purpose of protecting the low scale theory from large radiative corrections we impose supersymmetry. Specifically we investigate the constraints placed on models incorporating TeV scale supersymmetry.

A robust consequence of approximate left-right symmetry in the early Universe is the occurrence of transitory domain walls. It has been proposed [22] that such domain walls are susceptible to instability arising from nonrenormalizable operators suppressed by Planck scale. Supergravity then introduces two interesting ingredients not considered in the previous treatments of domain walls. First, the structure of the nonrenormalizable terms is dictated by the supersymmetry formalism[23–25]. On the other hand, one has to contend with the danger of gravitino overabundance [26,27]. In this paper we explore the restrictions on the possible energy scale of parity breaking imposed by these considerations.

We study these effects in the context of two implementations of left-right symmetry, one where all superfields carry nontrivial gauge couplings and another, for comparison, which contains a gauge singlet. We also study these models within three different scenarios for the dynamics of the wall complex. One is a scenario in which the walls disappear within the radiation dominated era, another where dominance of moduli keeps the Universe matter dominated during the domain wall evolution, and a third wherein the domain walls in fact come to dominate the Universe for a limited epoch, accompanied by a mild inflationary phase. In all the scenarios of domain wall evolution, the left-right symmetric model with a singlet turns out to be less restricted than the one without singlets. The model without any singlet turns out, at least in one scenario of wall evolution, to be sufficiently restrictive that the parity breaking scale can be no larger than 10^8 GeV. The overall lesson is that the new features introduced by supergravity can have a strong bearing on the scale of parity breaking for ensuring viable cosmology free of permanent domain walls.

The alternative to Planck scale suppressed terms for distinguishing between parity symmetric vacua was studied in [28,29] wherein the parity breaking operators are induced at a much lower scale, viz., the supersymmetry

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breaking scale, and signaled by the gauge mediation mechanism, thus linking the two scales. This avenue for evading unwanted domain walls remains open for models that get restricted in the present study.

In the remainder of the paper, in Sec. II we review cosmology with domain walls. In Sec. III we discuss three possible scenarios for evolution of domain walls and the requirement in each case for the successful destabilization of the wall complex. In Sec. IV we discuss the origin of the parity breaking terms within supergravity formalism. In Sec. V we discuss the essentials of supersymmetric left-right model. In subsection V A we discuss a particular renormalizable implementation of left-right supersymmetric model and then check for the sufficiency of the induced Planck scale terms to cause the required destabilization of the domain walls. The same study is carried out for a recent implementation of supersymmetric left-right symmetry including a gauge singlet in subsection V B. In Sec. VI we summarize the conclusions. Two appendixes A and B contain the calculations of the Planck scale suppressed terms for the two left-right symmetric models studied.

II. DISCRETE SYMMETRY BREAKING AND COSMOLOGY

Spontaneous breakdown of discrete symmetries in unified models can give rise to domain walls. Stable domain walls from unified theories have long been recognized as inconsistent with the observed Universe [30]. In the presence of several degenerate ground states the domain wall network in the Universe can be rather complicated. If the domain walls are stable, the energy density stored in the network decreases as $\rho \propto 1/a$, resulting in $a(t) \propto t^2$ leading to a mild inflationary behavior. A more generic possibility is that the walls continue to be destroyed due to collisions and result in the formation of homogeneous domains. However even one domain wall of grand unified scale of the size of the horizon can conflict with cosmology. For domain walls arising at symmetry breaking scale M_R, it can be estimated [30,31] that the density perturbation introduced by them would conflict with the known magnitude of temperature fluctuation $\Delta T/T \sim 10^{-5}$ of the cosmic microwave background if $M_R \gtrsim 1$ MeV. This impasse is overcome if the spontaneously broken discrete symmetry is also broken explicitly by a small amount. For example [32], the symmetry $\phi \to -\phi$ can be broken by adding a term $\epsilon\phi^3$ to the Lagrangian which gives a pressure difference governed by the small parameter ϵ between the two sides of the domain walls.

The authors of [22] have discussed several similar reasons for considering such gravity induced terms and their effect in destabilizing domain walls. For the theory of a generic neutral scalar field ϕ, the effective higher dimensional operators can be written as

$$V_{\text{eff}} = \frac{C_5}{M_{\text{Pl}}}\phi^5 + \frac{C_6}{M_{\text{Pl}}^2}\phi^6 + \cdots. \tag{1}$$

Such terms give rise to a pressure difference across a given domain wall of the amount of the difference in the effective energy density across the wall, $\delta\rho = \Delta V_{\text{eff}}$. Specifically, the terms odd in ϕ break the discrete symmetry, and the leading contribution to the difference in pressure is $\frac{2C_5}{M_{\text{Pl}}} \times \langle\phi\rangle^5$. From cosmological considerations we can separately estimate the difference $\delta\rho$ in the energy density across a domain wall required for timely removal of the domain walls. It is found that this has a value smaller by a factor M_R/M_{Pl} than the leading order term in the generic effective potential considered above, leading to the conclusion that the walls will indeed be removed without conflicting with cosmology.

This toy example however is only instructional because in realistic theories, gauge invariance and supersymmetry significantly constrain the structure of terms that can arise. In a nonsupersymmetric example [33], gauge invariance implies that the leading order operator is suppressed by $1/M_{\text{Pl}}^2$. Further, the terms are products of vacuum expectation values of different scalar fields which may differ significantly in their mass scales, as will be true in our study. The exceptional case where the toy example may be of direct relevance is the presence of one or more gauge singlet scalar fields in the theory (equivalently, superfields in a supersymmetric theory) permitting the kind of terms listed above.

As an illustration of this phenomenon consider the leading order operator containing several scalar fields, ϕ_i

$$\frac{\phi_1\phi_2\phi_3\phi_4\phi_5}{M_{\text{Pl}}}. \tag{2}$$

Borrowing from a calculation that will be detailed later [see Eq. (13)], suppose the required constraint for successful removal of domain walls is $\delta\rho \gtrsim M_R^{11/2}/M_{\text{Pl}}^{3/2}$. Then the requirement that the operator in Eq. (2) is sufficient for removing domain walls is that

$$\frac{\phi_1\phi_2\phi_3\phi_4\phi_5}{M_{\text{Pl}}} \gtrsim \frac{M_R^{11/2}}{M_{\text{Pl}}^{3/2}}. \tag{3}$$

Now suppose that there are only two kinds of fields, one getting the vacuum expectation value of the order of M_R, the parity breaking scale, and that there are x factors of this field in the operator, while the other field constitutes the remaining factors, and gets a TeV scale value v. Then

$$v^{5-x} \gtrsim \frac{M_R^{(11/2)-x}}{M_{\text{Pl}}^{1/2}}, \tag{4}$$

so that, taking TeV scale to be $v \sim 10^3$ GeV,

$$\log M_R \lesssim \frac{24.5 - 3x}{5.5 - x}. \tag{5}$$

This relation means that if $x = 5$, M_R can be as large as the Planck scale, while for $x = 1$, M_R is forced to have a value $<10^5$ GeV, In our analysis we shall be assuming that parity breaking occurs at the same mass scale as the mass of the heavy Majorana neutrinos. The constraints to be derived also depend on a few ancillary details, specifically the dynamics of the walls before they disintegrate. The primary implication of these other details, to be discussed in the following, is only the value of the temperature at which the standard cosmology resumes. In the following we shall admit the possibility of this more general kind of evolution and focus attention on two issues. The Universe should be radiation dominated at temperature 10 MeV and lower in order to ensure successful big-bang nucleosynthesis (BBN). Secondly, the danger of gravitino overabundance is generic to all supersymmetric models. Detailed calculations [26,34] show that the gravitinos with unacceptable consequences to observable cosmology are generated entirely after reheating of the Universe subsequent to primordial inflation provided $T_R \gtrsim 10^9$ GeV. Thus we make the requirement that entropy generated from the decay of domain walls should not raise the temperature of the Universe above this temperature scale.

III. MODELS OF DYNAMICS OF DOMAIN WALL COMPLEX

Occurrence of domain walls *per se* at some epoch in the early Universe is not ruled out, provided the walls eventually get destroyed. Safe disappearance of domain walls was dealt with in some generality in [22,35], the former in the context of Planck scale effects, and the latter in the context of instanton induced effects from QCD. There have been several model specific studies of the fate of domain walls, e.g., [36–38] and studies pointing out that transitory domain walls may in fact form the basis for explanation of some of the cosmological effects such as leptogenesis [33,38–40] or address problems such as proliferation of relics [28,37,38].

For the purpose of this paper we consider three possible routes through which domain walls may evolve. The first one consists of domain walls originating in a radiation dominated era and destabilized and destroyed also within the radiation dominated era before they begin to dominate the energy density of the Universe (referred to as the RD model). The second scenario was essentially proposed in [37], which consists of the walls originating in a radiation dominated phase, subsequent to which the Universe enters a "matter dominated" phase, either due to substantial production of heavy unwanted relics such as moduli, or due to the presence of a coherently oscillating scalar field (referred to as the MD model). Here too walls disappear before they come to dominate the energy density of the Universe. Finally, we consider a variant of the MD model in which the domain walls come to dominate the energy density of the Universe and continue to do so for a considerable epoch, leading to mild inflationary behavior or weak inflation [41,42] (referred to as the WI model). We now describe these in detail.

A. Evolution during a radiation dominated era

The essentials of this scenario are as originally proposed by Kibble [30] and Vilenkin [32]. Domain walls arise at some temperature T_c, the critical temperature of a phase transition at which a scalar field ϕ acquires a nonzero vacuum expectation value at a scale M_R. The energy density trapped per unit area of the wall is $\sigma \sim M_R^3$. The dynamics of the walls is determined by two quantities, force due to tension f_T and force due to friction f_F. The first of these is determined by intrinsic energy per unit area σ, and the average scale of radius of curvature R prevailing in the wall complex. We estimate $f_T \sim \sigma/R$. The frictional force is proportional to the collisions encountered by the wall with surrounding radiation with energy density $\sim T^4$, while the former is navigating through the medium at speed β. This force is estimated as $f_F \sim \beta T^4$. The epoch at which these two forces balance each other sets the time scale $t_R \sim R/\beta$. We may take this as the time scale by which the wall portions that started with radius of curvature scale R straighten out. Putting together these statements we get the following scaling law for the growth of the scale $R(t)$ on which the wall complex is smoothed out:

$$R(t) \approx (G\sigma)^{1/2} t^{3/2}. \tag{6}$$

Now the energy density of the domain walls goes as $\rho_W \sim (\sigma R^2/R^3) \sim (\sigma/Gt^3)^{1/2}$. In a radiation dominated era this ρ_W becomes comparable to the energy density of the Universe [$\rho \sim 1/(Gt^2)$] around time $t_0 \sim 1/(G\sigma)$.

Next, we consider destabilization of walls due to the pressure difference $\delta\rho$ arising from small asymmetry in the conditions on the two sides. This effect competes with the two quantities mentioned above. Since $f_F \sim 1/(Gt^2)$ and $f_T \sim (\sigma/(Gt^3))^{1/2}$, it is clear that at some point of time, $\delta\rho$ would exceed either the force due to tension or the force due to friction. For either of these requirements to be satisfied before $t_0 \sim 1/(G\sigma)$, we get

$$\delta\rho \geq G\sigma^2 \approx \frac{M_R^6}{M_{\rm Pl}^2} \sim M_R^4 \frac{M_R^2}{M_{\rm Pl}^2}. \tag{7}$$

We may read this formula by defining a factor

$$\mathcal{F} \equiv \frac{\delta\rho}{M_R^4}, \tag{8}$$

where M_R^4 is the energy density in the wall complex immediately at the phase transition, which relaxes by factor $\mathcal{F}$ at the epoch of its decay. The factor $\mathcal{F}$ is strongly dependent on the assumed model of evolution of the wall complex, and is found to be $M_R^2/M_{\rm Pl}^2$ in this model.

B. Evolution during a matter dominated era

We next take up the model of evolution in which the scale factor behaves as in a matter dominated era by the time the domain walls get destabilized. This possibility was considered in [37] by Kawasaki and Takahashi. The analysis begins by assuming that the initially formed wall complex in a phase transition is expected to rapidly relax to a few walls per horizon volume at an epoch characterized by Hubble parameter value H_i. Thus the initial energy density in the wall complex is $\rho_W^{(\mathrm{in})} \sim \sigma H_i$. This epoch onward, we assume the energy density of the Universe to be dominated by heavy relics or an oscillating modulus field, in either of which cases, the scale factor $a(t) \propto t^{2/3}$. The energy density for both of these cases scales as $\rho_{\mathrm{mod}} \sim \rho_{\mathrm{mod}}^{(\mathrm{in})}/a(t)^3$. If the domain wall (DW) complex remains frustrated, i.e. its energy density contribution $\rho_{\mathrm{DW}} \propto 1/a(t)$, it can be seen that [37] the Hubble parameter at the epoch of equality of DW contribution with contribution of the rest of the matter is given by

$$H_{\mathrm{eq}} \sim \sigma^{3/4} H_i^{1/4} M_{\mathrm{Pl}}^{-(3/2)}. \tag{9}$$

To proceed we assume that the domain walls start decaying as soon as they dominate the energy density of the Universe. If the temperature at this particular epoch is T_D, then $H_{\mathrm{eq}}^2 \sim G T_D^4$. So from Eq. (9) we find

$$T_D^4 \sim \sigma^{3/2} H_i^{1/2} M_{\mathrm{Pl}}^{-1}. \tag{10}$$

Let us assume the temperature at which the domain walls are formed $T \sim \sigma^{1/3}$. So

$$H_i^2 = \frac{8\pi}{3} G \sigma^{4/3} \sim \frac{\sigma^{4/3}}{M_{\mathrm{Pl}}^2}. \tag{11}$$

From Eq. (10), we get

$$T_D^4 \sim \frac{\sigma^{11/6}}{M_{\mathrm{Pl}}^{3/2}} \sim \frac{M_R^{11/2}}{M_{\mathrm{Pl}}^{3/2}} \sim M_R^4 \left(\frac{M_R}{M_{\mathrm{Pl}}}\right)^{3/2}. \tag{12}$$

Now requiring $\delta\rho > T_D^4$, we get

$$\delta\rho > M_R^4 \left(\frac{M_R}{M_{\mathrm{Pl}}}\right)^{3/2}. \tag{13}$$

Thus in this case we find $\mathcal{F} \equiv (M_R/M_{\mathrm{Pl}})^{3/2}$, a milder suppression factor than in the radiation dominated case above.

C. Evolution including weak inflation

The third possibility we consider is that the walls do not disintegrate by the time they come to dominate the energy density of the Universe, but in fact go on to dominate the energy density of the Universe. This domination however lasts for a limited epoch. Since the Universe evolves as $a(t) \propto t^2$, it leads to an epoch of mild inflation (as against exponential) also referred to as thermal or weak inflation. This possibility has been considered [41,42] in the context of removal moduli in superstring cosmology [43–45]. Such a situation is most likely in the case when the $\delta\rho$ is typically small, not large enough to destabilize the walls sufficiently quickly. But eventually a small pressure difference will also win over f_T or f_F, because either the curvature scale R diverges, as for straightened out walls, or the translational speed β reduces drastically. Since we have no microscopic model for deciding which of these is finally responsible, we introduce a temperature scale T_D at which the walls begin to experience instability. Note that unlike in the previous example, we will not be able to estimate T_D in terms of other mass scales and will accept it as undetermined and consider a few reasonable values for it for our final estimate.

As has been studied above, at H_{eq} the energy density of the domain wall network dominates energy density of the Universe. The scale factor at this epoch is characterized by a_{eq}. Denoting the energy density of the domain walls at the time of equality as $\rho_{\mathrm{DW}}(t_{\mathrm{eq}})$, the evolution of energy density can be written as

$$\rho_{\mathrm{DW}}(t_d) \sim \rho_{\mathrm{DW}}(t_{\mathrm{eq}}) \left(\frac{a_{\mathrm{eq}}}{a_d}\right), \tag{14}$$

where a_d is scale factor at the epoch of decay of domain walls corresponding to time t_d. If the domain walls decay at an epoch characterized by temperature T_D, then $\rho_{\mathrm{DW}}(t_d) \sim T_D^4$. So from Eq. (14),

$$T_D^4 = \rho_{\mathrm{DW}}(t_{\mathrm{eq}}) \left(\frac{a_{\mathrm{eq}}}{a_d}\right). \tag{15}$$

In the matter dominated era the energy density of the moduli fields scale as

$$\rho_{\mathrm{mod}}^d \sim \rho_{\mathrm{mod}}^{\mathrm{eq}} \left(\frac{a_{\mathrm{eq}}}{a_d}\right)^3. \tag{16}$$

Substituting the value of a_{eq}/a_d from Eq. (15) in the above equation,

$$\rho_{\mathrm{mod}}^d \sim \frac{T_D^{12}}{\rho_{\mathrm{DW}}^2(t_{\mathrm{eq}})}. \tag{17}$$

The energy density of the domain walls dominates this model universe after the time of equality, $\rho_{\mathrm{DW}}(t_d) > \rho_{\mathrm{mod}}^d$. So the pressure difference across the domain walls when they start decaying is given by

$$\delta\rho \gtrsim \frac{T_D^{12} G^2}{H_{\mathrm{eq}}^4}, \tag{18}$$

where we have used the relation $H_{\mathrm{eq}}^2 \sim G\rho_{\mathrm{DW}}(t_{\mathrm{eq}})$. Replacing the value of H_{eq} from Eq. (9), and $H_i^2 \sim G\rho_{\mathrm{DW}}^{\mathrm{in}} \sim M_R^4/M_{\mathrm{Pl}}^2$,

$$\delta\rho \gtrsim M_R^4 \left(\frac{T_D^{12} M_{\mathrm{Pl}}^3}{M_R^{15}}\right). \tag{19}$$

The $\mathcal{F}$ factor introduced in Eq. (8) turns out in this case to be $(T_D^{12} M_{\rm Pl}^3)/M_R^{15}$, rather sensitively dependent upon T_D.

IV. SUPERGRAVITY AND LEFT-RIGHT SYMMETRY BREAKING

The possibility that left-right symmetry may remain unbroken to low scales, and such breaking may be compatible with standard cosmology has been studied in an earlier work [29]. Specifically it was examined whether in the supersymmetric left-right symmetric model the parity breaking could be of hidden sector origin, and communicated to the visible sector through gauge mediation at a low scale. The attempt is to see if several of the puzzles of the standard model and incorporation of right-handed neutrinos are essentially possible within a few orders of magnitude of the TeV scale. It was found that messengers of a particular implementation of gauge mediated supersymmetry breaking can also communicate left-right symmetry violation. This is independent of the mechanism for the left-right symmetry violation in the hidden sector, the origin of which therefore remains unknown.

The question is, does supergravity have the potential to address the origin of left-right symmetry violation at a low scale? There is a folk theorem that discrete symmetries can be violated by quantum gravity effects. The reasoning runs as follows. Formation of black hole horizons can cause unaccounted violation of a global charge, while preserving gauge charges. We then expect black-holelike virtual states in quantum gravity which can induce effective terms violating the global charges. Such induced terms however do not arise in the process of perturbative renormalization, since every perturbative term, even in a nonrenormalizable Lagrangian should obey the expected symmetries. The effective terms would therefore arise from instantonlike effects.

Because of discrete nature of the symmetry, the signal of its breaking would be in the difference in the coefficients of the terms that get interchanged under the symmetry operation. The structure of supergravity ensures that at the renormalizable level gravity couples separately to the left sector and right sector with no mixing terms. The field contents are identical and the gauge couplings in the two sectors are identical at this order. It appears very difficult to see how supergravity would distinguish between the constants induced in the two sectors. We also note that $N = 1$ supergravity is finite at one-loop level. This leads us to suspect that we should not expect parity violating terms from supergravity, at least in the leading order in $1/M_{\rm Pl}$.

Thus a justification for considering $1/M_{\rm Pl}$ terms differing in their coefficients arises from the possibility that such are a result of gravity mediated supersymmetry breaking communicated from the hidden sector. For this to work we must assume one of two possibilities, either that the gauge group governing the hidden sector does not admit left-right symmetry as a subgroup or that such symmetry is broken in the hidden sector. The breaking should then be communicated to the visible sector along with the supersymmetry breaking. The root cause of the parity breaking then would remain hidden as in our earlier work. It should be emphasized that if this is the case, the terms signaling the breaking of parity must be proportional to the scale of supersymmetry breaking in the hidden sector and vanish in the limit of exact supersymmetry. In the formalism we have adopted we choose to order the terms by powers of inverse $M_{\rm Pl}$, however the induced terms must be understood to also be dependent on the scale of supersymmetry breaking.

The question of direct breaking of parity within supergravity is much less clear. However in this phenomenological analysis we shall also consider the next-to-leading order terms that would arise assuming that leading-order terms somehow do not break parity. According to Eq. (43) it is found that such terms may suffice only marginally to solve the problem of unacceptable domain walls in cosmology.

Regardless of their origin, the structure of the symmetry breaking terms in the scalar potential will be similar to that of the terms that can be derived from the superpotential, as happens in the case of soft supersymmetry breaking terms [23]. Thus we may use the usual supergravity formalism to derive the terms through which Planck suppressed left-right symmetry breaking may get manifested. A similar approach has been adopted in the context of the minimal supersymmetric standard model (MSSM) in [46] where the origin and the effect of higher dimensional operators have been discussed in the context of collider data. In the remainder of this section we summarize the essential formalism to be used in our calculation. We adopt the notation described in [23]. The supergravity Lagrangian is obtained from three functions of complex scalar fields, viz., superpotential (W), Kähler potential (K), and gauge kinetic function f_{ab}. The F-term contribution to the scalar potential in supergravity theory

$$V_F = k_i^j F_j F^{*i} - 3e^{K/M_{\rm Pl}^2} W W^*/M_{\rm Pl}^2, \tag{20}$$

where

$$F^i = -[(K^{-1})^i_l (W^{*l} + W^* K^l/M_{\rm Pl}^2)]. \tag{21}$$

Making use of Eq. (21) in Eq. (20) the individual terms in V_F can be written as

$$V_F = (K^{-1})^{*k}_l \bigg[W_k^* W^l + \frac{W_k^* W K^l}{M_{\rm Pl}^2} + \frac{W^* K_k W^l}{M_{\rm Pl}^2} + \frac{W^* K_k W K^l}{M_{\rm Pl}^4} \bigg] + \text{higher order terms}. \tag{22}$$

Here we have considered the first term of Eq. (20). The scalar potential contains D-term contributions from gauge interactions which are given by

$$V_D = \frac{1}{2}\,\mathrm{Re}[f_{ab}^{-1}\hat{D}_a\hat{D}_b], \tag{23}$$

where

$$\hat{D}_a = -K^i(T^a)_i{}^j\phi_j = -\phi^{*j}(T^a)_j{}^iK_i \tag{24}$$

and f^{ab} is the gauge kinetic function which is given by

$$f^{ab} = \delta_{ab}[1/g_a^2 + f_a^i\phi_i/M_{\rm Pl} + \cdots]. \tag{25}$$

For our purpose it will be sufficient to consider $f^{ab} = \delta_{ab}/g_a^2$ since we do not expect left-right asymmetry to arise from the gauge sector. In the following we shall consider the terms arising in the scalar potential from expanding W and K in the powers of $(1/M_{\rm Pl})$.

V. GRAVITY INDUCED OPERATORS IN LEFT-RIGHT SUPERSYMMETRIC MODEL

The key difference in any realistic model in comparison with the generic considerations of [22] is that gauge invariance restricts the structure of the nonrenormalizable terms. For instance in [33] the operators considered were

$$V_{\text{non-SUSY}} \sim c_1\frac{1}{M_{\rm Pl}^2}(\Delta_L^\dagger\Delta_L)^3 + c_2\frac{1}{M_{\rm Pl}^2}(\Delta_R^\dagger\Delta_R)^3. \tag{26}$$

In other words, gauge invariance requires the terms to be $O(1/M_{\rm Pl}^2)$ rather than $O(1/M_{\rm Pl})$. A difference in pressure caused by such operators, after putting vacuum expectation values of the fields $\Delta \sim M_R$, would be adequate to remove the domain walls accompanied by radiation dominated era evolution, Eq. (7), and even more so the domain walls that evolved during an effective matter dominated era, Eq. (13).

But in the models we consider, supersymmetry forbids the kind of terms shown in Eq. (26); on the other hand, generic parity breaking terms of $O(1/M_{\rm Pl})$ respecting $SU(2)_L$ and $SU(2)_R$ gauge invariance are permitted. This is due to additional field content with different gauge charges. This gain in order of $1/M_{\rm Pl}$ is however offset by the fact that, due to demands of phenomenology, some of the fields can acquire TeV scale vacuum expectation values as well. Here we have studied two successful supersymmetric implementations of left-right symmetry but the method can be extended to other implementations.

The minimal supersymmetric left-right model is based on the gauge group $SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$. The anomaly free $B-L$ global symmetry of the standard model is promoted to a gauge symmetry. The quark, lepton, and Higgs superfields for the minimal supersymmetric left-right model, with their respective quantum numbers under the gauge group $SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ are given by

$$\begin{aligned} &Q = (3,2,1,1/3), && Q_c = (3^*,1,2,-1/3),\\ &L = (1,2,1,-1), && L_c = (1,1,2,1),\\ &\Phi_i = (1,2,2,0), && \text{for } i = 1,2,\\ &\Delta = (1,3,1,2), && \bar{\Delta} = (1,3,1,-2),\\ &\Delta_c = (1,1,3,-2), && \bar{\Delta}_c = (1,1,3,2), \end{aligned} \tag{27}$$

where we have suppressed the generation index for simplicity of notation. In the Higgs sector, the bidoublet Φ is doubled to have nonvanishing Cabbibo-Kobayashi-Maskawa matrix, whereas the Δ triplets are doubled to have anomaly cancellation. Under discrete parity symmetry the fields are prescribed to transform as

$$\begin{gathered} Q \leftrightarrow Q_c^*, \qquad L \leftrightarrow L_c^*, \qquad \Phi_i \leftrightarrow \Phi_i^\dagger,\\ \Delta \leftrightarrow \Delta_c^*, \qquad \bar{\Delta} \leftrightarrow \bar{\Delta}_c^*. \end{gathered} \tag{28}$$

However, this minimal left-right symmetric model is unable to break parity spontaneously [14,15]. Inclusion of nonrenormalizable terms gives a more realistic structure of possible vacua [16,17,47]. Such terms were studied for the case when the scale of $SU(2)_R$ breaking is high, close to Planck scale.

A. The ABMRS model with a pair of triplets

Because of the difficulties with the model discussed above, a "minimal" renormalizable model was developed by Aulakh *et al.* early in [16–18] and will be referred to here as the ABMRS model after the authors. In this model two triplet fields, Ω and Ω_c, were added, with the following quantum numbers:

$$\Omega = (1,3,1,0), \qquad \Omega_c = (1,1,3,0), \tag{29}$$

which was shown to improve the situation with only the renormalizable terms [17,18,48]. It was shown that this model breaks down to the minimal supersymmetric standard model at low scale. This model was studied in the context of cosmology in [28,38] and, specifically, the mechanism for leptogenesis via domain walls in [49].

The superpotential for this model including Higgs fields only is given by

$$\begin{aligned} W_{LR} = {}& m_\Delta\,\mathrm{Tr}\Delta\bar{\Delta} + m_\Delta\,\mathrm{Tr}\Delta_c\bar{\Delta}_c + \frac{m_\Omega}{2}\mathrm{Tr}\Omega^2\\ &+ \frac{m_\Omega}{2}\mathrm{Tr}\Omega_c^2 + \mu_{ij}\,\mathrm{Tr}\tau_2\Phi_i^T\tau_2\Phi_j + a\,\mathrm{Tr}\Delta\Omega\bar{\Delta}\\ &+ a\,\mathrm{Tr}\Delta_c\Omega_c\bar{\Delta}_c + \alpha_{ij}\,\mathrm{Tr}\Omega\Phi_i\tau_2\Phi_j^T\tau_2\\ &+ \alpha_{ij}\,\mathrm{Tr}\Omega_c\Phi_i^T\tau_2\Phi_j\tau_2. \end{aligned} \tag{30}$$

Since supersymmetry is broken at a very low scale, we can employ the F and D flatness conditions obtained from the superpotential to get a possible solution for the vacuum expectation values (VEV's for the Higgs fields,

$$\langle\Omega\rangle = 0, \quad \langle\Delta\rangle = 0, \quad \langle\bar{\Delta}\rangle = 0,$$
$$\langle\Omega_c\rangle = \begin{pmatrix} \omega_c & 0 \\ 0 & -\omega_c \end{pmatrix}, \tag{31}$$
$$\langle\Delta_c\rangle = \begin{pmatrix} 0 & 0 \\ d_c & 0 \end{pmatrix}, \quad \langle\bar{\Delta}_c\rangle = \begin{pmatrix} 0 & \bar{d}_c \\ 0 & 0 \end{pmatrix}.$$

This solution set is of course not unique. Since the original theory is parity invariant a second solution for the F and D flat conditions exists, with left-type fields' VEV's exchanged with those of the right-type fields [38,49].

With VEV's as in Eq. (31) the pattern of breaking is

$$SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L} \xrightarrow{M_R} SU(2)_L \otimes U(1)_R \otimes U(1)_{B-L}, \tag{32}$$

$$\xrightarrow{M_{B-L}} SU(2)_L \otimes U(1)_Y. \tag{33}$$

The model introduces a new mass scale, m_Ω. However, it was observed [17] that these terms can be forbidden in the superpotential by invoking an R symmetry, and then the corresponding terms appearing in the scalar potential can be interpreted as soft terms entering only after supersymmetry breakdown at the electroweak scale. This approach imposes a condition on the scales of breaking, with respect to the electroweak scale M_W,

$$M_R M_W \simeq M_{B-L}^2. \tag{34}$$

This relation raises the interesting possibility that the scale of M_R can be as low as 10^4 to 10^6 GeV, with corresponding very low scale 10^3 to 10^4 GeV of lepton number violation, opening the possibility of low energy leptogenesis [19,49].

As discussed in Sec. (IV) we shall proceed to find the $1/M_{\rm Pl}$ terms in the effective potential by expanding Kähler potential and superpotential in powers of $1/M_{\rm Pl}$. Here we include $\Delta(\bar{\Delta})$, $\Delta_c(\bar{\Delta}_c)$, and $\Omega(\Omega_c)$ fields in the expansion of the Kähler potential and superpotential. The Kähler potential in this model up to $(1/M_{\rm Pl})$ can be written as

$$K = {\rm Tr}(\Delta\Delta^\dagger + \bar{\Delta}\bar{\Delta}^\dagger) + {\rm Tr}(\Delta_c\Delta_c^\dagger + \bar{\Delta}_c\bar{\Delta}_c^\dagger) + {\rm Tr}(\Omega\Omega^\dagger) + {\rm Tr}(\Omega_c\Omega_c^\dagger) + \frac{C_L}{M_{\rm Pl}}({\rm Tr}\Delta\Omega\Delta^\dagger + {\rm Tr}\bar{\Delta}\Omega\bar{\Delta}^\dagger + {\rm H.c.}) + \frac{C_R}{M_{\rm Pl}}({\rm Tr}\Delta_c\Omega_c\Delta_c^\dagger + {\rm Tr}\bar{\Delta}_c\Omega_c\bar{\Delta}_c^\dagger + {\rm H.c.}), \tag{35}$$

where C_L and C_R are dimensionless constants. The superpotential can also be expanded in powers of $(1/M_{\rm Pl})$ which is written as

$$\begin{aligned} W &= W_{\rm ren} + W_{\rm nr} \\ &= m_\Delta({\rm Tr}\Delta\bar{\Delta} + {\rm Tr}\Delta_c\bar{\Delta}_c) + \frac{m_\Omega}{2}({\rm Tr}\Omega^2 + {\rm Tr}\Omega_c^2) + a({\rm Tr}\Delta\Omega\bar{\Delta} + {\rm Tr}\Delta_c\Omega_c\bar{\Delta}_c) + \frac{a_L}{2M_{\rm Pl}}({\rm Tr}\Delta\bar{\Delta})^2 + \frac{a_R}{2M_{\rm Pl}}({\rm Tr}\Delta_c\bar{\Delta}_c)^2 \\ &\quad + \frac{b_L}{M_{\rm Pl}}{\rm Tr}\Delta^2\,{\rm Tr}\bar{\Delta}^2 + \frac{b_R}{M_{\rm Pl}}{\rm Tr}\Delta_c^2\,{\rm Tr}\bar{\Delta}_c^2 + \frac{c_L}{4M_{\rm Pl}}({\rm Tr}\Omega^2)^2 + \frac{c_R}{4M_{\rm Pl}}({\rm Tr}\Omega_c^2)^2 + \frac{d_L}{2M_{\rm Pl}}{\rm Tr}\Omega^2\,{\rm Tr}\Delta\bar{\Delta} + \frac{d_R}{2M_{\rm Pl}}{\rm Tr}\Omega_c^2\,{\rm Tr}\Delta_c\bar{\Delta}_c \\ &\quad + \frac{f}{M_{\rm Pl}}{\rm Tr}\Delta\bar{\Delta}\,{\rm Tr}\Delta_c\bar{\Delta}_c + \frac{h_1}{M_{\rm Pl}}{\rm Tr}\Delta^2\,{\rm Tr}\Delta_c^2 + \frac{h_2}{M_{\rm Pl}}{\rm Tr}\bar{\Delta}^2\,{\rm Tr}\bar{\Delta}_c^2 + \frac{j}{M_{\rm Pl}}{\rm Tr}\Omega^2\,{\rm Tr}\Omega_c^2 + \frac{k}{M_{\rm Pl}}{\rm Tr}\Omega^2\,{\rm Tr}\Delta_c\bar{\Delta}_c \\ &\quad + \frac{m}{M_{\rm Pl}}{\rm Tr}\Omega_c^2\,{\rm Tr}\Delta\bar{\Delta}, \end{aligned} \tag{36}$$

where the coefficients appearing against the individual terms are dimensionless constants.

The effective potential resulting from the above modifications is calculated in Appendix A. We have also examined the D terms in the effective potential and find that they do not give rise to $O(\frac{1}{M_{\rm Pl}})$ terms. Thus we only need to deal with F terms. Assuming a phase in which the right-type fields get a nontrivial VEV and all left-type fields have zero VEV, the expression for the leading term in $1/M_{\rm Pl}$, the scalar potential becomes

$$V_{\rm eff}^R \sim \frac{a(c_R + d_R)}{M_{\rm Pl}} M_R^4 M_W + \frac{a(a_R + d_R)}{M_{\rm Pl}} M_R^3 M_W^2. \tag{37}$$

Now due to left-right symmetry, there is also a corresponding phase in which the left-type fields get a VEV, but not the right-type fields. For this phase the value of the effective potential is

$$V_{\rm eff}^L \sim \frac{a(c_L + d_L)}{M_{\rm Pl}} M_R^4 M_W + \frac{a(a_L + d_L)}{M_{\rm Pl}} M_R^3 M_W^2. \tag{38}$$

The possibility of these two phases of approximately equal energy density gives rise to domain walls separating such phases. The pressure difference across the walls is proportional to the difference in energy density between two sides of the wall, and is given by

$$\begin{aligned} \delta\rho &\sim [(c_L - c_R) + (d_L - d_R)]\frac{M_R^4 M_W}{M_{\rm Pl}} + [(a_L - a_R) + (d_L - d_R)]\frac{M_R^3 M_W^2}{M_{\rm Pl}} \\ &\sim \kappa^A \frac{M_R^4 M_W}{M_{\rm Pl}} + \kappa'^A \frac{M_R^3 M_W^2}{M_{\rm Pl}}, \end{aligned} \tag{39}$$

where $\kappa^A = (c_L - c_R) + (d_L - d_R)$ and $\kappa'^A = (a_L - a_R) + (d_L - d_R)$, and the superscript A refers to the ABMRS model. From Eq. (39), we see that to leading order in $1/M_{\rm Pl}$ there are two kinds of operators appearing in $\delta\rho$, differing in powers of (M_W/M_R).

We shall now compare these operators with the energy density required for the successful removal of the domain walls in the three cases labeled as RD, MD, and WI, respectively, discussed in Sec. II. Comparing Eq. (7) with individual operators in Eq. (39) and taking the scale M_R as 10^6 GeV, and taking the more dominant term κ, we get the constraint

$$\kappa^A_{\rm RD} > 10^{-10}\left(\frac{M_R}{10^6\ {\rm GeV}}\right)^2. \tag{40}$$

This is easily satisfied at the low scale of M_R proposed. For M_R scale tuned to 10^9 GeV, needed to avoid the gravitino problem after reheating at the end of inflation, $\kappa_{\rm RD} \sim 10^{-4}$, is a reasonable constraint. but requires $\kappa^A_{\rm RD}$ to be $O(1)$ if the scale of M_R is an intermediate scale 10^{11} GeV.

Next, comparing Eq. (13) with individual terms in Eq. (39), the constraint on $\kappa^A_{\rm MD}$ is found to be

$$\kappa^A_{\rm MD} > 10^{-2}\left(\frac{M_R}{10^6\ {\rm GeV}}\right)^{3/2}, \tag{41}$$

which puts a modest requirement on the value of $\kappa^A_{\rm MD}$ for suitable disappearance of domain walls. However taking $M_R \sim 10^9$ GeV being the temperature scale required to have thermal leptogenesis, without the undesirable gravitino production, leads to $\kappa^A_{\rm MD} > 10^{5/2}$, an unacceptable requirement. The MD case is in fact generic to supersymmetric and string inspired models [43–45] due to moduli production. And we find that in this case the ABMRS model requires a low scale of M_R and nonthermal or resonant leptogenesis.

In the WI case, Eq. (19) there is extreme sensitivity to the scales of M_R and T_D. Proceeding in same way as above comparing Eq. (19) with Eq. (39) the constraint on κ is found to be

$$\kappa^A_{\rm WI} > 10^{-4}\left(\frac{10^6\ {\rm GeV}}{M_R}\right)^{15}\left(\frac{T_D}{10\ {\rm GeV}}\right)^{12}. \tag{42}$$

This is a reasonable constraint for the proposed median values of the two mass scales. However the constraint makes the model rather strongly predictive. The scale of decay of the wall complex T_D can be any value below the chosen M_R scale. Thus if $T_D \sim 10^4$ GeV, then M_R is forced to be closer to the gravitino scale 10^9 GeV. This can be problematic if the reheating temperature after the disappearance of the domain walls is comparable to the temperature scale of the original phase transition. The Universe would reheat to 10^9 GeV, raising the possibility of unacceptable gravitino regeneration.

Finally we consider the possibility, raised in Sec. IV, that the gravity induced terms are of direct origin, and due to one-loop finiteness of supergravity, do not give rise to $1/M_{\rm Pl}$ terms in the superpotential or the Kähler potential. In such a case the most dominant operator to be considered is suppressed by $(1/M_{\rm Pl})^2$. In the ABMRS case we find, after substituting the vacuum expectation values, that such an operator has the magnitude

$$\delta V_{\text{next-to-leading order}} \sim \frac{M_R^4 M_W^2}{M_{\rm Pl}^2}. \tag{43}$$

So long as we are considering theories with M_R values less than an intermediate scale 10^{11} GeV, such terms are subdominant to the ones considered above. However if the leading terms are absent, the constraint on the coefficient for the above term for each of the above constraints is tightened by a factor $M_{\rm Pl}/M_W \sim 10^{16}$. Such a constraint immediately renders the first two scenarios of domain wall evolution cosmologically unacceptable. The third case of weak inflation however continues to be possible for phenomenologically acceptable values of the energy scales.

To summarize the situation for the ABMRS model, we have found that there is an upper bound on the scale M_R if the wall evolution unfolds during a radiation dominated epoch or a matter dominated epoch. In the latter case, which is generic for string theory cosmology with the presence of heavy moduli fields, the natural value of M_R is required to be significantly lower than 10^9 GeV. In the case of an evolution accompanied by a weak inflationary epoch, there is no upper bound, rather a lower bound on the scale M_R but which is extremely sensitive to the value of the scale T_D at which the walls may finally disappear.

B. The BM model with a single singlet

An independent approach to improve the minimal model with introduction of a parity odd singlet [50] was adopted in [14,15]. However this was shown at tree level to lead to charge-breaking vacua being at a lower potential than charge-preserving vacua.

Recently, an alternative to this has been considered in Babu and Mohapatra in [51] where a superfield $S = (1, 1, 1, 0)$, also singlet under parity, is included in addition to the minimal set of Higgs required, as in Eq. (27). The Δ_c, $\bar{\Delta}_c$ are required for $SU(2)_R \otimes U(1)_{B-L}$ symmetry breaking without inducing R-parity violating couplings. The singlet field S is introduced so that $SU(2)_R \otimes U(1)_{B-L}$ symmetry breaking occurs in the supersymmetric limit. We refer to this as the BM model, after its authors. The superpotential is given by

$$W_{LR} = W^{(1)} + W^{(2)},$$

where

$$\begin{aligned} W^{(1)} = {} & \mathbf{h}_l^{(i)} L^T \tau_2 \Phi_i \tau_2 L_c + \mathbf{h}_q^{(i)} Q^T \tau_2 \Phi_i \tau_2 Q_c \\ & + i\mathbf{f}^* L^T \tau_2 \Delta L + i\mathbf{f} L^{cT} \tau_2 \Delta_c L_c + S[\lambda^*\ {\rm Tr}\Delta\bar{\Delta} \\ & + \lambda\,{\rm Tr}\Delta_c\bar{\Delta}_c + \lambda'_{ab}\,{\rm Tr}\Phi_a^T \tau_2 \Phi_b \tau_2 - M_R^2], \end{aligned} \tag{44}$$

$$W^{(2)} = M_\Delta \operatorname{Tr}\Delta\bar{\Delta} + M_\Delta^* \operatorname{Tr}\Delta_c\bar{\Delta}_c + \mu_{ab}\operatorname{Tr}\Phi_a^T\tau_2\Phi_b\tau_2 + M_s S^2 + \lambda_s S^3. \quad (45)$$

In this analysis the terms in $W^{(2)}$ have been assumed to be zero. Dropping the terms in $W^{(2)}$ makes the theory more symmetric and more predictive. It is observed that dropping quadratic and cubic terms in S leads to an enhanced R-symmetry. Further, dropping the massive couplings introduced for Δ's means that Δ masses arise purely from supersymmetry (SUSY) breaking effects, keeping these fields light and relevant to collider phenomenology. Dropping the μ_{ab} terms for Φ fields makes it possible to explain the μ parameter of MSSM as being spontaneously induced from S VEV through terms in $W^{(1)}$. Additionally, absence of the $W^{(2)}$ terms can be shown to solve the SUSY CP and strong CP problems.

The presence of linear terms in S in $W^{(1)}$ makes possible the following SUSY vacuum:

$$\langle S\rangle = 0, \qquad \lambda v_R\bar{v}_R + \lambda^* v_L\bar{v}_L = M_R^2, \quad (46)$$

where $v_L(\bar{v}_L)$ and $v_R(\bar{v}_R)$ are the VEV's of the neutral components of $\Delta(\bar{\Delta})$ and $\Delta_c(\bar{\Delta}_c)$ fields, respectively. From Eq. (46) it is clear that we have a flat direction in the $v_L - v_R$ space. Assuming that the flat directions are lifted, we have two choices, viz.,

$$v_R = \bar{v}_R = 0, \qquad |v_L| = |\bar{v}_L| = \frac{M_R}{\sqrt{\lambda^*}}, \quad (47)$$

$$v_L = \bar{v}_L = 0, \qquad |v_R| = |\bar{v}_R| = \frac{M_R}{\sqrt{\lambda}}. \quad (48)$$

The important result is that after SUSY breaking and emergence of SUSY breaking soft terms, integrating out heavy sleptons modifies the vacuum structure due to Coleman-Weinberg type one-loop terms which must be treated to be of the same order as the other terms in V^{eff}. Accordingly, it is shown [51] that the V^{eff} contains terms of the form

$$V^{\mathrm{eff}}_{\text{one-loop}}(\Delta_c) \sim -|f|^2 m_{L^c}^2 \operatorname{Tr}(\Delta_c\Delta_c^\dagger)A_1^R - |f|^2 m_{L^c}^2 \operatorname{Tr}(\Delta_c\Delta_c)\operatorname{Tr}(\Delta_c^\dagger\Delta_c^\dagger)A_2^R, \quad (49)$$

where A_1^R and A_2^R are constants obtained from expansion of the effective potential. Presence of these terms is shown to lead to the consequence of a preference for the electric charge-preserving vacuum over the charge-breaking vacuum, provided $m_{L^c}^2 < 0$.

Further Eq. (47) also constitutes a valid solution of Eq. (46). In this vacuum the soft terms can give rise to the following terms in the effective potential:

$$V^{\mathrm{eff}}_{\text{one-loop}}(\Delta) \sim -|f|^2 m_L^2 \operatorname{Tr}(\Delta\Delta^\dagger)A_1^L - |f|^2 m_L^2 \operatorname{Tr}(\Delta\Delta)\operatorname{Tr}(\Delta^\dagger\Delta^\dagger)A_2^L, \quad (50)$$

with A_1^L and A_2^L constants. Thus the choice of known phenomenology is only one of two possible local choices, and formation of domain walls is inevitable.

Here we analyze the full superpotential without setting $W^{(2)}$ terms to zero. As in the ABMRS model, we study here the scalar potential taking supergravity into account. Including triplet fields and singlet field S, the Kähler potential in this model up to $O(1/M_{\mathrm{Pl}})$ is given by

$$\begin{aligned} K = {} & \operatorname{Tr}(\Delta\Delta^\dagger + \bar{\Delta}\bar{\Delta}^\dagger) + \operatorname{Tr}(\Delta_c\Delta_c^\dagger + \bar{\Delta}_c\bar{\Delta}_c^\dagger) + |S|^2 \\ & + \frac{C_L}{M_{\mathrm{Pl}}}(\operatorname{Tr}\Delta S\Delta^\dagger + \operatorname{Tr}\bar{\Delta}S\bar{\Delta}^\dagger + \text{H.c.}) \\ & + \frac{C_R}{M_{\mathrm{Pl}}}(\operatorname{Tr}\Delta_c S\Delta_c^\dagger + \operatorname{Tr}\bar{\Delta}_c S\bar{\Delta}_c^\dagger + \text{H.c.}) \\ & + \frac{d}{M_{\mathrm{Pl}}}(S^3 + \text{H.c.}). \end{aligned} \quad (51)$$

The superpotential upt o $(1/M_{\mathrm{Pl}})$ order is given by

$$\begin{aligned} W = {} & W_{\mathrm{ren}} + W_{\mathrm{nr}} \\ = {} & m_\Delta(\operatorname{Tr}\Delta\bar{\Delta} + \operatorname{Tr}\Delta_c\bar{\Delta}_c) + M_s S^2 + \lambda_s S^3 \\ & + S[\lambda^*\operatorname{Tr}\Delta\bar{\Delta} + \lambda\operatorname{Tr}\Delta_c\bar{\Delta}_c - M_R^2] + \frac{a_L}{2M_{\mathrm{Pl}}}(\operatorname{Tr}\Delta\bar{\Delta})^2 \\ & + \frac{a_R}{2M_{\mathrm{Pl}}}(\operatorname{Tr}\Delta_c\bar{\Delta}_c)^2 + \frac{b_L}{2M_{\mathrm{Pl}}}\operatorname{Tr}\Delta^2\operatorname{Tr}\bar{\Delta}^2 \\ & + \frac{b_R}{2M_{\mathrm{Pl}}}\operatorname{Tr}\Delta_c^2\operatorname{Tr}\bar{\Delta}_c^2 + \frac{c}{M_{\mathrm{Pl}}}S^4 + \frac{c_L}{M_{\mathrm{Pl}}}S^2\operatorname{Tr}\Delta\bar{\Delta} \\ & + \frac{c_R}{M_{\mathrm{Pl}}}S^2\operatorname{Tr}\Delta_c\bar{\Delta}_c + \frac{f}{M_{\mathrm{Pl}}}\operatorname{Tr}\Delta\bar{\Delta}\operatorname{Tr}\Delta_c\bar{\Delta}_c \\ & + \frac{h_1}{M_{\mathrm{Pl}}}\operatorname{Tr}\Delta^2\operatorname{Tr}\Delta_c^2 + \frac{h_2}{M_{\mathrm{Pl}}}\operatorname{Tr}\bar{\Delta}^2\operatorname{Tr}\bar{\Delta}_c^2. \end{aligned} \quad (52)$$

The effective potential has been calculated in Appendix B considering the term from the effective potential $(K^{-1})^{*k}{}_l W_k^* W^l$. When the right-type fields get a VEV the scalar potential can be written as

$$V^R_{\mathrm{eff}} \sim \frac{a_R}{M_{\mathrm{Pl}}}M_R^5 + \frac{a_R}{M_{\mathrm{Pl}}}sM_R^4 + \frac{C_R}{M_{\mathrm{Pl}}}s^2M_R^3 + \frac{C_R}{M_{\mathrm{Pl}}}s^3M_R^2, \quad (53)$$

where s is the scale at which S gets a VEV. To calculate the potential when only the left-type fields get vacuum expectation values, we introduce corresponding coefficients a_L, etc. We then compute the pressure difference across the walls as

$$\begin{aligned} \delta\rho & \sim (a_L - a_R)\frac{M_R^5}{M_{\mathrm{Pl}}} + \cdots + (C_L - C_R)\frac{s^3M_R^2}{M_{\mathrm{Pl}}} \\ & \sim \kappa^B\frac{M_R^5}{M_{\mathrm{Pl}}} + \cdots + \kappa'^B\frac{s^3M_R^2}{M_{\mathrm{Pl}}}, \end{aligned} \quad (54)$$

where $\kappa^B = (a_L - a_R)$, $\kappa'^B = (C_L - C_R)$, superscript B referring to the BM model and the ellipses $(\ldots)$ are in lieu of terms which, as we explain next, are relatively unimportant. According to the BM model, the value s is of the scale of supersymmetry breaking. If this scale is TeV scale, then M_R is expected to be higher and the κ'^B terms are subdominant. In case however the supersymmetry breaking scale is 10^{11} GeV, it could be higher than the scale of M_R. In this case the κ'^B term is expected to dominate.

In the case of TeV scale supersymmetry breaking, comparing the first term in Eq. (54) with Eqs. (7), (13), and (19) of with Sec. II, we get the corresponding constraints on possible values of κ as

$$\kappa^B_{\mathrm{RD}} > 10^{-13}\left(\frac{M_R}{10^6\ \mathrm{GeV}}\right), \tag{55}$$

$$\kappa^B_{\mathrm{MD}} > 10^{-6}\left(\frac{M_R}{10^6\ \mathrm{GeV}}\right)^{1/2}, \tag{56}$$

$$\kappa^B_{\mathrm{WI}} > 10^{-8}\left(\frac{10^6\ \mathrm{GeV}}{M_R}\right)^{16}\left(\frac{T_D}{10\ \mathrm{GeV}}\right)^{12}. \tag{57}$$

Thus for the proposed M_R scale of 10^6 GeV there is no serious constraint on κ^B values. Only in the scenario with weak inflation, if the T_D scale is high, such as 100 GeV, the value $M_R \sim 10^6$ GeV becomes marginal, but due to large powers of the mass scale present, a small increase in M_R easily offsets the effect of the increase in T_D. Overall, the kind of operators obtained in this particular model provides no constraint on the mass scale M_R, as long as the scale of supersymmetry breaking is TeV scale.

To check other possibilities, we consider the supersymmetry breaking scale M_S and hence s to be $\sim 10^{11}$ GeV. In this case we get, proceeding as above,

$$\kappa'^B_{\mathrm{RD}} > 10^{-25}\left(\frac{M_R}{10^6\ \mathrm{GeV}}\right)^4\left(\frac{10^{11}\ \mathrm{GeV}}{s}\right)^3, \tag{58}$$

$$\kappa'^B_{\mathrm{MD}} > 10^{-19}\left(\frac{M_R}{10^6\ \mathrm{GeV}}\right)^{7/2}\left(\frac{10^{11}\ \mathrm{GeV}}{s}\right)^3, \tag{59}$$

$$\kappa'^B_{\mathrm{WI}} > 10^{-44}\left(\frac{T_D}{10\ \mathrm{GeV}}\right)^{12}\left(\frac{10^6\ \mathrm{GeV}}{M_R}\right)^{17}\left(\frac{10^{11}\ \mathrm{GeV}}{s}\right)^3. \tag{60}$$

This shows that there is no particular constraint on the induced parity breaking coefficients due to an increase in the scale of supersymmetry breaking M_S, so long as $M_R < s \equiv M_S \sim 10^{11}$ GeV.

In summary, the BM model remains mostly unrestricted by the present considerations. This is due to newer terms possible with a gauge singlet.

VI. CONCLUSION

Ever since the discovery of massive neutrinos, it has become a tantalizing possibility that the small neutrino masses arise from rich physics at a high energy scale, which in turn would naturally incorporate right-handed neutrinos. A model that treats this new content symmetrically with the known contents would naturally lead to the requirement of parity symmetry in the high energy model. If this discrete symmetry is spontaneously broken it would lead to formation of domain walls in the early Universe.

We have considered three scenarios for the evolution of transitory domain wall networks ending in their decay. We characterize each model by a dimensionless parameter $\mathcal{F}$ which is the ratio of the available pressure difference across a wall at the time of its decay to the characteristic energy density M_R^4 in the Universe at the time of formation of the wall complex; see Eqs. (7), (13), and (19). Of the three scenarios, the first one unfolds entirely in a radiation dominated universe, in which the dynamics is governed by the interplay of forces due to friction and tension and the pressure difference across the walls. Here we find that the parameter $\mathcal{F}$ is given by $(M_R/M_{\mathrm{Pl}})^2$. The second scenario unfolds in a matter dominated era, where the domain wall complex decays as soon as the energy density of the same dominates the energy density of the Universe. The parameter $\mathcal{F}$ in this case is given by $(M_R/M_{\mathrm{Pl}})^{3/2}$. The third scenario is an extension of the second one where we assume that the domain wall network comes to dominate the energy density of the Universe and continues to do so for a finite epoch before it decays. Characterizing the required pressure difference $\delta\rho$ across the domain walls in this case requires additional input, the ratio of scale factor values a_{eq}/a_d where subscript eq refers to the epoch at which the domain walls become equally as important as the rest of the matter and subscript d refers to the epoch at which the decay of the domain walls occurs. We characterize this ratio by an equivalent "decay temperature" T_D defined in Eq. (15). The formula for $\mathcal{F}$ in this case shows very high sensitivity to the mass scales concerned.

For each of these cases we study the viability of the two specific models of spontaneous parity breaking, the ABMRS model in Sec. V A and the BM model in Sec. V B. Each of the particle physics models permits intrinsic operators whose coefficients must match up to the required parameter $\mathcal{F}$, resulting in final destabilization of the wall complex. If the operators available in the model cannot provide a $\delta\rho$ of required magnitude, the wall complex would not be destabilized, leading to unacceptable cosmology.

The ABMRS model turns out to be more restrictive, using as it does only nontrivial representations of the gauge group. In this case a high scale for the parity breaking becomes conditionally disfavored, though still viable if the wall evolution leads to a weak inflationary epoch. The BM model containing a singlet turns out to not be restricted by the considerations here.

Our main conclusion is that a low scale scenario with $M_R \sim 10^6$ GeV or lower is viable and generic. Specifically, in the ABMRS model with domain wall evolution in a matter dominated epoch M_R is restricted to remain less than 10^8 GeV; Eq. (41). A matter dominated epoch is generic to string theory inspired models with the occurrence of moduli fields of mass scale 10^9 GeV and hence this restriction is of special interest.

The mechanism studied here is an alternative to an earlier one [29] that studied the possibility of parity breaking mediated by the messengers in a version of the gauge mediated supersymmetry breaking scenario. The constraints on the parameters for that scenario to ensure disappearance of domain walls were rather stringent. We now see that wall disappearance by Planck scale suppressed terms is more plausible. The precise source of parity breaking is not identified in our analysis, but assumed to occur in a hidden sector. Our parity breaking terms are ordered as a series in $1/M_{\rm Pl}$, but the terms would also be proportional to the scale of supersymmetry breaking, and would vanish in the limit of exact supersymmetry. Our broad conclusion is that Planck scale suppressed terms can suffice for removal of unwanted domain walls in realistic theories with broken supersymmetry, but can imply important constraints on the concerned energy scales.

ACKNOWLEDGMENTS

U. A. Y. records his thanks to Guy Moore for critical comments.

APPENDIX A: ABMRS MODEL

The scalar potential contains D-term contributions from gauge interactions. From Eq. (23)

$$\begin{aligned} V_D &= \frac{1}{2}\,\mathrm{Re}(g_a^2/\delta_{ab}\hat{D}_a\hat{D}_b) \\ &= \frac{g_a^2}{2}\,\mathrm{Re}(\hat{D}_a\hat{D}_a) \\ &= \frac{g_a^2}{2}\,\mathrm{Re}\,\mathrm{Tr}(K_\Delta(T^a)\Delta) + \mathrm{Tr}(K_{\bar\Delta}(T^a)\bar\Delta) + \mathrm{Tr}(K_\Omega(T^a)\Omega)^2 \\ &= \frac{g^2}{8}\,\mathrm{Re}\,\mathrm{Tr}(K_\Delta\tau^a\Delta) + \mathrm{Tr}(K_{\bar\Delta}\tau^a\bar\Delta) + \mathrm{Tr}(K_\Omega\tau^a\Omega)^2 + \frac{g'^2}{8}\,\mathrm{Re}\,\mathrm{Tr}(K_\Delta\Delta) - \mathrm{Tr}(K_{\bar\Delta}\bar\Delta)^2 \\ &= \frac{g^2}{8}\,\mathrm{Re}2\,\mathrm{Tr}(\Delta^\dagger\tau^a\Delta) + \frac{2C_L}{M_{\rm Pl}}\mathrm{Tr}(\Omega\Delta^\dagger\tau^a\Delta) + 2\,\mathrm{Tr}(\bar\Delta^\dagger\tau^a\bar\Delta) + \frac{2C_L}{M_{\rm Pl}}\mathrm{Tr}(\Omega\bar\Delta^\dagger\tau^a\bar\Delta) + 4\,\mathrm{Tr}(\Omega\tau^a\Omega) \\ &\quad + \frac{2C_L}{M_{\rm Pl}}[\mathrm{Tr}(\Delta^\dagger\Delta\tau^a\Omega) + \mathrm{Tr}(\bar\Delta^\dagger\bar\Delta\tau^a\Omega)]^2 + \frac{g'^2}{8}\,\mathrm{Re}2\,\mathrm{Tr}(\Delta^\dagger\Delta) - 2\,\mathrm{Tr}(\bar\Delta^\dagger\bar\Delta) + \frac{2C_L}{M_{\rm Pl}}\mathrm{Tr}(\Delta\Omega\Delta^\dagger) - \frac{2C_L}{M_{\rm Pl}}\mathrm{Tr}(\bar\Delta\Omega\bar\Delta^\dagger)^2. \end{aligned} \tag{A1}$$

The D term vanishes after putting the VEV's for the corresponding fields. From above it is clear that we cannot find $1/M_{\rm Pl}$ suppressed terms from V_D. So we have to go for V_F to find the desired terms. Here we consider the first term appearing in Eq. (22), i.e. $(k^{-1})^{*k}_l W^*_k W^l$.

Substituting the Eqs. (36) and (35) in Eq. (22), the terms which contribute are

$$\begin{aligned} V_R \sim{} & \frac{ad_R}{2M_{\rm Pl}}\mathrm{Tr}\Omega_c^2\,\mathrm{Tr}\Omega_c\bar\Delta_c\bar\Delta_c^\dagger + \frac{a_R}{M_{\rm Pl}}m_\Delta\,\mathrm{Tr}\Delta_c\bar\Delta_c\,\mathrm{Tr}\bar\Delta_c\bar\Delta_c^\dagger \\ & + \frac{a_R}{M_{\rm Pl}}m_\Delta\,\mathrm{Tr}\Delta_c\bar\Delta_c\,\mathrm{Tr}\Delta_c\Delta_c^\dagger + \frac{ad_R}{2M_{\rm Pl}}\mathrm{Tr}\Omega_c^2\,\mathrm{Tr}\Omega_c\Delta_c\Delta_c^\dagger \\ & + \frac{c_R}{M_{\rm Pl}}m_\Omega(\mathrm{Tr}\Omega_c^2)^2 + \frac{ac_R}{M_{\rm Pl}}\mathrm{Tr}\Omega_c^2\,\mathrm{Tr}\Delta_c\Omega_c\bar\Delta_c \\ & + \text{terms higher order in } 1/M_{\rm Pl}. \end{aligned} \tag{A2}$$

In the ABMRS model we have the relation

$$\begin{aligned} M_{B-L}^2 \simeq M_R M_W; \qquad \omega = -\left|\frac{m_\Delta}{a}\right| \equiv M_R; \\ d = \bar d = (2m_\Delta m_\Omega/a^2)^{1/2} \equiv M_{B-L}. \end{aligned} \tag{A3}$$

After putting the VEV's for the corresponding fields and making use of appropriate scale, the terms up to $O(1/M_{\rm Pl})$ are

$$V_R \sim \frac{a(c_R + d_R)}{M_{\rm Pl}}M_R^4 M_W + \frac{a(a_R + d_R)}{M_{\rm Pl}}M_R^3 M_W^2. \tag{A4}$$

APPENDIX B: BM MODEL

The leading order terms in this model come from the first term of Eq. (22). Writing explicitly the individual terms,

$$\begin{aligned} V_F ={} & (K^{-1})_{\Delta_c\Delta_c^\dagger}W^*_{\Delta_c}W^{\Delta_c} + (K^{-1})_{\bar\Delta_c\bar\Delta_c^\dagger}W^*_{\bar\Delta_c}W^{\bar\Delta_c} \\ & + (K^{-1})_{SS^*}W^*_S W^S. \end{aligned} \tag{B1}$$

Calculating the above terms, the terms in the scalar potential in lowest order in $1/M_{\rm Pl}$ are given by

$$V_F \sim \frac{a_R}{M_{\rm Pl}} m_\Delta \,{\rm Tr}\Delta_c \bar{\Delta}_c \,{\rm Tr}\bar{\Delta}_c \bar{\Delta}_c^\dagger + \frac{a_R}{M_{\rm Pl}} \lambda^* S \,{\rm Tr}\Delta_c \bar{\Delta}_c \,{\rm Tr}\bar{\Delta}_c \bar{\Delta}_c^\dagger + \frac{C_R}{M_{\rm Pl}} [\lambda^2 S^3 \,{\rm Tr}\Delta_c \Delta_c^\dagger + M_\Delta^2 S \,{\rm Tr}\Delta_c \Delta_c^\dagger + S^2(\lambda^* M_\Delta \,{\rm Tr}\Delta_c \Delta_c^\dagger + \lambda M_\Delta \,{\rm Tr}\Delta_c \Delta_c^\dagger)] + \frac{C_R}{M_{\rm Pl}} \lambda S \,{\rm Tr}\Delta_c \bar{\Delta}_c \,{\rm Tr}\bar{\Delta}_c \bar{\Delta}_c^\dagger + \frac{C_R}{M_{\rm Pl}} M_s S^2 \,{\rm Tr}\Delta_c \bar{\Delta}_c + \frac{C_R}{M_{\rm Pl}} s^3 \,{\rm Tr}\Delta_c \bar{\Delta}_c + \text{other terms.} \tag{B2}$$

After putting the VEV's for the neutral components of the triplet field and using the appropriate scale, the term in the highest power of M_R

$$V_R \sim \frac{a_R}{M_{\rm Pl}} M_R^5 + \frac{a_R}{M_{\rm Pl}} s M_R^4 + \frac{C_R}{M_{\rm Pl}} s^2 M_R^3 + \frac{C_R}{M_{\rm Pl}} s^3 M_R^2 + \text{other terms.} \tag{B3}$$

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