

Gauge mediated supersymmetry breaking and the cosmology of the left-right symmetric model

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Left-right symmetry including supersymmetry presents an important class of gauge models which may possess natural solutions to many issues of phenomenology. Cosmology of such models indicates a phase transition accompanied by domain walls. Such walls must be unstable in order to not conflict with standard cosmology and can further be shown to assist with open issues of cosmology such as dilution of unwanted relic densities and leptogenesis. In this paper we construct a model of gauge mediated supersymmetry breaking in which parity breaking is also signaled along with supersymmetry breaking and so as to be consistent with cosmological requirements. It is shown that addressing all the stated cosmological issues requires an extent of fine-tuning, while in the absence of fine-tuning, leptogenesis accompanying successful completion of the phase transition is still viable.

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I. INTRODUCTION

The left-right symmetric model [1–5] has received considerable attention as a simple extension of the standard model (SM). It provides an elegant explanation of several open questions of the SM in addition to providing a natural explanation for the smallness of neutrino masses [6–9] via the seesaw mechanism [10–13]. In [14,15] the possibility that it be consistent as a TeV scale extension of the SM was explored. In the supersymmetric setting it was shown by [16] that $M_R \sim \text{TeV}$ is consistent with $SO(10)$ at $M_X \sim 10^{16}$ GeV. In this paper we adopt a similar approach viz., the scale of left-right symmetry can be at a TeV to PeV scale, and that supersymmetry¹ (SUSY) is preserved down to the same scale, so that the TeV/PeV scale is protected from unreasonable corrections from higher energy scales [20,21].

The cosmological phase transition accompanying the parity breakdown is effectively a first order phase transition wherein domains of two different kinds of vacua separated by a network of domain walls (DW) occur [22]. The need for ensuring a homogeneous universe in late cosmology requires the phase transition to end appropriately, in particular, that this network of domain walls be unstable. The limits on the epoch to which the DW network may survive are placed by the requirements of successful big bang nucleosynthesis and cosmic microwave background radiation data. A possible signature for a phase transition ending

with decay of domain walls at such energy scales is relic gravitational waves detectable at upcoming experiments [23].

The unstable domain walls can have specific physical consequences. One of the possibilities is that they are very slow in disappearing. This is generically true if the difference in values of the effective potential across the wall is small. The energy density of the domain wall complex scales with the cosmological scale factor a as $\rho \propto 1/a$, resulting in $a(t) \propto t^2$ leading to a mild inflationary behavior. This is in accord with a proposal of thermal inflation [24] (in some contexts called weak inflation) which can help to dilute unwanted relics arising in string theory [25,26].

Unstable domain walls also provide the nonadiabatic conditions for leptogenesis. Theories with Majorana neutrino masses and especially with gauged $B - L$ symmetry present the interesting possibility of explaining the baryon asymmetry of the Universe via thermal processes in the early Universe [27]. This particular approach however has been shown to generically require the scale of Majorana neutrino mass, equivalently, the scale of gauged $B - L$ symmetry breaking in relevant models, to be $10^{11} - 10^{13}$ GeV [28,29], with a more optimistic constraint $M_{B-L} > 10^9$ GeV [30,31]. On the other hand, it has been shown [32,33] that the only real requirement imposed by leptogenesis is that the presence of heavy neutrinos should not erase lepton asymmetry generated by a given mechanism, possibly nonthermal. This places the modest bound $M_1 > 10^4$ GeV, on the mass of the lightest of the heavy Majorana neutrinos.

There are two possible scenarios which exploit this window of low mass scale for leptogenesis. One is the “soft leptogenesis” [34–37], relying on the decay of scalar

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¹For reviews of supersymmetry, see [17–19] and references therein.

superpartners of neutrino and a high degree of degeneracy [38] in the mass eigenvalues due to soft SUSY breaking terms. Another possibility for leptogenesis is provided by the unstable domain walls. This mechanism is analogous to that explored for the electroweak baryogenesis [39], provided a source for CP asymmetry can be found. It has been shown [40] that the domain walls occurring in models such as are studied in this paper generically give spatially varying complex masses to neutrinos. Unstable domain walls of our models are therefore sufficient to ensure the required leptogenesis.

In an earlier work two of the authors explored the possibility of consistent cosmology in this class of models, primarily the questions of removal of unwanted relics [41] and the role of domain walls as possible catalyzers of leptogenesis [42]. Exact left-right symmetry of the underlying model however does not permit instability of the domain walls. In [41] it was shown that this circumstance is avoided provided supersymmetry breaking in the hidden sector also breaks parity symmetry and this is signaled through the soft terms.

Several proposals have been made for successful phenomenological implementation of SUSY in a left-right symmetric gauge theory [43–47]. Implications of gauge mediation of SUSY in the same was studied in [48,49]. More recently, this model has been studied in the context of the CERN LHC in [50,51] and also a similar low energy model in the context of cosmology in [52]. In all such models considered, spontaneous gauge symmetry breaking required to recover SM phenomenology also leads to observed parity breaking. However, for cosmological reasons it is not sufficient to ensure local breakdown of parity. It has been proposed earlier [53] that the occurrence of the SM-like sector globally, i.e., homogeneously over the entire Universe, could be connected to the SUSY breaking effects from the hidden sector.

Here we pursue this question in the context of two classes of models, one represented by Aulakh, Benakli, Melfo, Rasin, and Senjanovic (ABMRS) [45,47] and the other, proposed recently by Babu and Mohapatra (BM) [51], both of which circumvent the thorny issue of phenomenologically acceptable vacua by making minimal extensions of the basic scheme. The former uses two triplets Ω and Ω_c neutral under $U(1)_{B-L}$ while the latter uses one superfield S singlet under all the proposed gauge interactions as well as parity. In the approach we adopt, namely, preservation of the left-right symmetry to low energies, both of these models have a generic issue regarding the occurrence of domain walls in cosmology which needs to be addressed. However the walls can be easily removed without conflicting with phenomenology even with very small parity breaking effects. In gauge mediated supersymmetry breaking (GMSB), the SUSY breaking effects are naturally small due to being communicated at one-loop and two-loop orders. We propose the required

(GMSB) scheme to obtain natural conditions for SUSY breaking and also for the disappearance of the domain walls. We show that the smallness migrates into parity breaking and is adequate to ensure global choice of true ground state in the Universe.

In Sec. II we discuss the two models ABMRS and BM. In Sec. III we discuss the soft terms arising in these models and the constraints imposed by consistent cosmology, presenting some of the concerned formulas in the Appendix. In Sec. IV we pursue the consequence of implementing generic GMSB in these theories. In Sec. V we propose a modification of the scheme to implement combined breaking of SUSY as well as parity in a manner consistent with low energy symmetries and successful cosmology. Section VI contains discussion of our results.

II. MINIMAL SUPERSYMMETRIC LEFT-RIGHT MODEL: A RECAP

The quark, lepton, and Higgs fields for the minimal left-right SUSY 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

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

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

$$\begin{aligned} Q &\leftrightarrow Q_c^*, & L &\leftrightarrow L_c^*, & \Phi_i &\leftrightarrow \Phi_i^\dagger, \\ \Delta &\leftrightarrow \Delta_c^*, & \bar{\Delta} &\leftrightarrow \bar{\Delta}_c^*. \end{aligned} \quad (2)$$

However, this minimal left-right symmetric model is unable to break parity spontaneously [43,44]. The inclusion of nonrenormalizable terms gives a more realistic structure of possible vacua [45,46,54]. Such terms were studied for the case when the scale of $SU(2)_R$ breaking is high, close to the Planck scale. We shall not pursue this possibility further here, retaining interest in TeV to PeV scale phenomenology.

A. The ABMRS model with a pair of triplets

Because of difficulties with the model discussed above, an early model to be called minimal by its authors is ABMRS [45–47]. Here two triplet fields Ω and Ω_c were added, with the following quantum numbers:

$$\Omega = (1, 3, 1, 0), \quad \Omega_c = (1, 1, 3, 0), \quad (3)$$

which was shown to improve the situation with only the renormalizable terms [45,47,55]. It was shown that this model breaks down to the minimal supersymmetric standard model (MSSM) at low scale. This model was studied in the context of cosmology in [41,53] specifically, the mechanism for leptogenesis via domain walls in [42].

The superpotential for this model is given by

$$\begin{aligned} W_{LR} = & \mathbf{h}_l^{(i)} L^T \tau_2 \Phi_i \tau_2 L_c + \mathbf{h}_q^{(i)} Q^T \tau_2 \Phi_i \tau_2 Q_c + i \mathbf{f} L^T \tau_2 \Delta L \\ & + i \mathbf{f} L^{cT} \tau_2 \Delta_c L_c + m_\Delta \text{Tr} \Delta \bar{\Delta} + m_\Delta \text{Tr} \Delta_c \bar{\Delta}_c \\ & + \frac{m_\Omega}{2} \text{Tr} \Omega^2 + \frac{m_{\Omega_c}}{2} \text{Tr} \Omega_c^2 + \mu_{ij} \text{Tr} \tau_2 \Phi_i^T \tau_2 \Phi_j \\ & + a \text{Tr} \Delta \Omega \bar{\Delta} + a \text{Tr} \Delta_c \Omega_c \bar{\Delta}_c + \alpha_{ij} \text{Tr} \Omega \Phi_i \tau_2 \Phi_j^T \tau_2 \\ & + \alpha_{ij} \text{Tr} \Omega_c \Phi_i^T \tau_2 \Phi_j \tau_2. \end{aligned} \quad (4)$$

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:

$$\begin{aligned} \langle \Omega \rangle = 0, \quad \langle \Delta \rangle = 0, \quad \langle \bar{\Delta} \rangle = 0, \quad \langle \Omega_c \rangle = \begin{pmatrix} \omega_c & 0 \\ 0 & -\omega_c \end{pmatrix}, \\ \langle \Delta_c \rangle = \begin{pmatrix} 0 & 0 \\ d_c & 0 \end{pmatrix}, \quad \langle \bar{\Delta}_c \rangle = \begin{pmatrix} 0 & \bar{d}_c \\ 0 & 0 \end{pmatrix}. \end{aligned} \quad (5)$$

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 [41,42].

With vev's as in Eq. (5) the pattern of breaking is

$$SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L} \xrightarrow{M_R} SU(2)_L \otimes U(1)_R \otimes U(1)_{B-L} \quad (6)$$

$$\xrightarrow{M_{B-L}} SU(2)_L \otimes U(1)_Y. \quad (7)$$

It was observed [45] that supersymmetric breaking imposes a condition on the scales of breaking, with respect to the electroweak scale M_W ,

$$M_R M_W \simeq M_{B-L}^2. \quad (8)$$

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 [32,42].

B. The BM model with a single singlet

An independent approach to improve the minimal model with the introduction of a parity odd singlet [56] was adopted in [43,44]. 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 [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. (1). The superpotential is given by

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

where

$$\begin{aligned} W^{(1)} = & \mathbf{h}_l^{(i)} L^T \tau_2 \Phi_i \tau_2 L_c + \mathbf{h}_q^{(i)} Q^T \tau_2 \Phi_i \tau_2 Q_c \\ & + i \mathbf{f}^* L^T \tau_2 \Delta L + i \mathbf{f} L^{cT} \tau_2 \Delta_c L_c + S[\lambda^* \text{Tr} \Delta \bar{\Delta} \\ & + \lambda \text{Tr} \Delta_c \bar{\Delta}_c + \lambda'_{ab} \text{Tr} \Phi_a^T \tau_2 \Phi_b \tau_2 - M_R^2], \end{aligned} \quad (9)$$

$$\begin{aligned} W^{(2)} = & M_\Delta \text{Tr} \Delta \bar{\Delta} + M_\Delta^* \text{Tr} \Delta_c \bar{\Delta}_c + \mu_{ab} \text{Tr} \Phi_a^T \tau_2 \Phi_b \tau_2 \\ & + M_s S^2 + \lambda_s S^3. \end{aligned} \quad (10)$$

For a variety of phenomenological reasons [51], the terms in $W^{(2)}$ may be 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 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, the 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:

$$\langle S \rangle = 0, \quad \lambda v_R \bar{v}_R + \lambda^* v_L \bar{v}_L = M_R^2, \quad (11)$$

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. In the ABMRS model, the introduction of a separate Ω field for each of the sectors L and R permits local preference of one sector over the other through spontaneous symmetry breaking. This preference however is local, valid in a finite sized region, giving rise to the possibility of domain walls. In the BM model however, due to the S field being neutral including under parity, such a distinction cannot arise even locally. This is reflected in the above equation (11) where we have a flat direction in the $v_L - v_R$ space.² A more general treatment of the possible vacua is included in the Appendix. Nevertheless,

$$v_L = \bar{v}_L = 0, \quad |v_R| = |\bar{v}_R| = \frac{M_R}{\sqrt{\lambda}} \quad (12)$$

is a possible vacuum [51] in which we recover the known phenomenology.

²The first equation of Sec. 3 of [51] is a special case of our condition (11).

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

$$V_{1\text{-loop}}^{\text{eff}}(\Delta_c) \sim -|f|^2 m_{L_c}^2 \text{Tr}(\Delta_c \Delta_c^\dagger) A_1^R - |f|^2 m_{L_c}^2 \text{Tr}(\Delta_c \Delta_c) \text{Tr}(\Delta_c^\dagger \Delta_c^\dagger) A_2^R, \quad (13)$$

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 happy consequence of a preference for the electric charge-preserving vacuum over the charge-breaking vacuum, provided $m_{L_c}^2 < 0$.

For the purpose of the present paper it is important to note that even assuming that some soft terms will lift the flat direction in (11), we still have no source of breaking $L - R$ symmetry. This means that

$$v_R = \bar{v}_R = 0, \quad |v_L| = |\bar{v}_L| = \frac{M_R}{\sqrt{\lambda^*}} \quad (14)$$

also constitutes a valid solution of Eq. (11). In this vacuum the soft terms can give rise to the following terms in the effective potential:

$$V_{1\text{-loop}}^{\text{eff}}(\Delta) \sim -|f|^2 m_L^2 \text{Tr}(\Delta \Delta^\dagger) A_1^L - |f|^2 m_L^2 \text{Tr}(\Delta \Delta) \text{Tr}(\Delta^\dagger \Delta^\dagger) A_2^L \quad (15)$$

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.

III. COSMOLOGY OF BREAKING AND SOFT TERMS

Because of the existence of two different sets of solutions for the possible vev's, formation of DWs is inevitable [41,42] in both the models considered above. Stable walls are known to overclose the Universe [22,57] and are undesirable. However, a small inequality in the free energies in the vacua on the two sides of the walls is sufficient to destabilize them. The lower bound on the temperature up to which such walls may persist in the Universe is set by the known physics of big bang nucleosynthesis, $T \sim 1$ MeV. Let the temperature by which the wall complex has substantially decayed, in particular, has ceased to dominate the energy density of the Universe, be T_D . It has been estimated that the free energy density difference $\delta\rho$ between the vacua which determines the pressure difference across a domain wall should be of the form [58,59]

$$\delta\rho \sim T_D^4 \quad (16)$$

in order for the DW to disappear at the scale T_D .

It has been observed in [60] that parity breaking effects suppressed by the Planck scale are sufficient to remove the DW. Black holes carry only locally conserved charges, and therefore processes in quantum gravity would exist which do not conserve global charges as well as violate discrete symmetry such as parity. Thus we expect parity breaking terms induced by quantum gravity in the effective Lagrangian. Assuming that the pressure across the walls is created by terms such as $C\varphi^6/M_{\text{Pl}}^2$, it is easy to check that a reasonable range of values of the order parameters $\langle\varphi\rangle$ and T_D exist for which the walls can disappear without conflicting with cosmology. This is especially true of high scale models, $M_R \gtrsim 10^{11}$ GeV. We shall be interested in a low (PeV) scale model where Planck scale effects can be ignored and where the parity breaking should arise from known effects which can be counterchecked against other observables.

In the ABMRS model at the scale M_R , $SU(2)_R \otimes U(1)_{B-L}$ breaks to $U(1)_R \otimes U(1)_{B-L}$, equally well, $SU(2)_L \otimes U(1)_{B-L}$ breaks to $U(1)_L \otimes U(1)_{B-L}$ depending on which of the two Ω fields acquires a vev. Thus domain walls are formed at the scale M_R . At a lower scale M_{B-L} when the Higgs triplet Δ 's get vev, $U(1)_R \otimes U(1)_{B-L}$ breaks to $U(1)_Y$ or $U(1)_L \otimes U(1)_{B-L}$ breaks to $U(1)_{Y'}$. In the BM model the breaking is directly to $SU(2)_L \otimes U(1)_Y$ or equally well, $SU(2)_R \otimes U(1)_{Y'}$. As per the analysis reported above we are assuming that the S field does not acquire a vev. Thus in each of the models we have MSSM after the M_{B-L} or the M_R scale, respectively. SUSY breaking soft terms emerge below the SUSY breaking scale M_S . We now proceed with the stipulation advanced in [53] in which the role of the hidden sector dynamics is not only to break SUSY but also break parity. This permits in principle a relation between observables arising from the two apparently independent breaking effects.

The soft terms which arise in the two models ABMRS and BM may be parametrized as follows:

$$\mathcal{L}_{\text{soft}}^1 = m_1^2 \text{Tr}(\Delta \Delta^\dagger) + m_2^2 \text{Tr}(\bar{\Delta} \bar{\Delta}^\dagger) + m_3^2 \text{Tr}(\Delta_c \Delta_c^\dagger) + m_4^2 \text{Tr}(\bar{\Delta}_c \bar{\Delta}_c^\dagger), \quad (17)$$

$$\mathcal{L}_{\text{soft}}^2 = \alpha_1 \text{Tr}(\Delta \Omega \Delta^\dagger) + \alpha_2 \text{Tr}(\bar{\Delta} \Omega \bar{\Delta}^\dagger) + \alpha_3 \text{Tr}(\Delta_c \Omega_c \Delta_c^\dagger) + \alpha_4 \text{Tr}(\bar{\Delta}_c \Omega_c \bar{\Delta}_c^\dagger), \quad (18)$$

$$\mathcal{L}_{\text{soft}}^3 = \beta_1 \text{Tr}(\Omega \Omega^\dagger) + \beta_2 \text{Tr}(\Omega_c \Omega_c^\dagger), \quad (19)$$

$$\mathcal{L}_{\text{soft}}^4 = S[\gamma_1 \text{Tr}(\Delta \Delta^\dagger) + \gamma_2 \text{Tr}(\bar{\Delta} \bar{\Delta}^\dagger)] + S^*[\gamma_3 \text{Tr}(\Delta_c \Delta_c^\dagger) + \gamma_4 \text{Tr}(\bar{\Delta}_c \bar{\Delta}_c^\dagger)], \quad (20)$$

$$\mathcal{L}_{\text{soft}}^5 = \tilde{\sigma}^2 |S|^2. \quad (21)$$

For the ABMRS model the relevant soft terms are given by

TABLE I. Differences in values of soft supersymmetry breaking parameters for a range of domain wall decay temperature values T_D . The differences signify the extent of parity breaking.

$T_D/\text{GeV}\sim$	10	10^2	10^3
$(m^2 - m'^2)/\text{GeV}^2\sim$	10^{-4}	1	10^4
$(\beta_1 - \beta_2)/\text{GeV}^2\sim$	10^{-8}	10^{-4}	1

$$\mathcal{L}_{\text{soft}} = \mathcal{L}_{\text{soft}}^1 + \mathcal{L}_{\text{soft}}^2 + \mathcal{L}_{\text{soft}}^3. \quad (22)$$

For the BM model the soft terms are given by

$$\mathcal{L}_{\text{soft}} = \mathcal{L}_{\text{soft}}^1 + \mathcal{L}_{\text{soft}}^4 + \mathcal{L}_{\text{soft}}^5. \quad (23)$$

Using the requirement of Eq. (16) we can constrain the differences between the soft terms in the left and right sectors [41,42]. According to Eq. (11) the S field does not acquire a vev in the physically relevant vacua and hence the terms in Eqs. (20) and (21) do not contribute to the vacuum energy. The terms in Eq. (18) are suppressed in magnitude relative to those in Eq. (19) due to having Ω vev's to one power lower. This argument assumes that the magnitude of the coefficients α are such as to not mix up the symmetry breaking scales of the Ω 's and the Δ 's.

To obtain orders of magnitude we have taken the m_i^2 parameters to be of the form $m_1^2 \sim m_2^2 \sim m^2$ and $m_3^2 \sim m_4^2 \sim m'^2$ [42] with T_D in the range $10\text{--}10^3$ GeV [26]. For both the models we have taken the value of the Δ vev's as $d \sim 10^4$ GeV. For the ABMRS model additionally we take $\omega \sim 10^6$ GeV. The resulting differences required for successful removal of domain walls are shown in Table I.

We see from Table I that assuming both the mass-squared differences $m^2 - m'^2$ and $\beta_1 - \beta_2$ arise from the same dynamics, Ω fields are the determinant of the cosmology. This is because the lower bound on the wall disappearance temperature T_D required by Ω fields is higher and the corresponding T_D is reached sooner. This situation changes if for some reason Ω 's do not contribute to the pressure difference across the walls. The BM model does not have Ω 's and falls in this category.

During the period of time in between destabilization of the DW and their decay, leptogenesis occurs due to these unstable DWs as discussed in [40,42]. After the disappearance of the walls at the scale T_D , electroweak symmetry breaks at a scale $M_{\text{EW}} \sim 10^2$ GeV and standard cosmology takes over. In the next section we discuss the implementation of the GMSB scenario for these models.

IV. GAUGE MEDIATED SUSY BREAKING IN MINIMAL SUPERSYMMETRIC MODELS

Proposals for an unobserved strongly coupled gauge sector giving rise to SUSY breaking which is then communicated to observable phenomenology were made in [61–67]. In proposals of gauge mediated SUSY breaking viable

at a low scale [68–70], dynamical breaking of SUSY in a hidden sector is communicated to a visible sector through a field X singlet under visible gauge interactions, and one or more conjugate pairs of chiral superfields called messenger fields which together constitute a vectorlike, anomaly free representation. For reviews the reader may refer to [17,71,72]. The choice of charges for the messenger fields ensures that they do not spoil gauge coupling unification. A simple choice is to choose them to be complete representations of the possible grand unification group, which in the left-right symmetric case is $SO(10)$.

The dynamical SUSY breakdown causes the messenger fields to develop SUSY violating interactions with the hidden sector, while they also interact with the (s)quarks, (s)leptons, and Higgs(inos) via gauge and gaugino interactions. Gaugino fields get a mass at one loop due to these interactions, while gauge invariance prevents gauge fields from acquiring any mass. As such supersymmetry is broken in the visible sector.

We denote the messenger sector fields to be Φ_i and its complex conjugate representation to be $\bar{\Phi}_i$, with i indexing the set of several possible messengers. These couple to a chiral superfield singlet X via a Yukawa-type interaction,

$$W = \sum_i y_i X \Phi_i \bar{\Phi}_i. \quad (24)$$

It should be noted that in the case of the BM model, this X could be identified with S . Coupling of X to the hidden sector gives rise to vacuum expectation values $\langle X \rangle$ and $\langle F_X \rangle$ for its scalar and auxiliary parts, respectively. As such the fermionic and scalar parts of the messenger sector get masses,

$$m_f^2 = |y_i \langle X \rangle|^2, \quad m_s^2 = |y_i \langle X \rangle|^2 \pm |y_i \langle F_X \rangle|. \quad (25)$$

Thus, the degeneracy between the fermionic and the scalar part of the messenger sector vanishes.

Gaugino mass arises due to one-loop diagrams and is given [69]

$$M_a = \frac{\alpha_a \langle F_X \rangle}{4\pi \langle X \rangle} (1 + O(x)), \quad (a = 1, 2, 3), \quad (26)$$

where $x = \langle F_X \rangle / (\lambda \langle X \rangle^2)$. Masses for the scalars of the SUSY model for the left-right symmetric case arise due to two-loop corrections and are given by

$$m_\phi^2 = 2 \left(\frac{\langle F_X \rangle}{\langle X \rangle} \right)^2 \left[\left(\frac{\alpha_3}{4\pi} \right)^2 C_3^\phi + \left(\frac{\alpha_2}{4\pi} \right)^2 (C_{2L}^\phi + C_{2R}^\phi) + \left(\frac{\alpha_1}{4\pi} \right)^2 C_1^\phi \right] (1 + O(x)). \quad (27)$$

The C_a^ϕ are the Casimir group theory invariants defined by

$$C_a^\phi \delta_i^j = (T^a T^a)_i^j, \quad (28)$$

where T^a is the group generator of the group which acts on the scalar ϕ . Since we consider both the models (ABMRS

and BM), the values of C_a^ϕ s for the fields are given by

$$\begin{aligned}
C_3^\phi &= \begin{cases} 4/3 & \text{for } \phi = Q, Q_c, \\ 0 & \text{for } \phi = \Phi_i, \Delta's, \Omega's, L's, S, \end{cases} \\
C_{2L}^\phi &= \begin{cases} 3/4 & \text{for } \phi = Q, L, \Phi_i, \\ 2 & \text{for } \phi = \Delta, \bar{\Delta}, \Omega, \\ 0 & \text{for } \phi = Q_c, L_c, \Delta_c, \bar{\Delta}_c, \Omega_c, S, \end{cases} \quad (29) \\
C_{2R}^\phi &= \begin{cases} 3/4 & \text{for } \phi = Q_c, L_c, \Phi_i, \\ 2 & \text{for } \phi = \Delta_c, \bar{\Delta}_c, \Omega_c, \\ 0 & \text{for } \phi = Q, L, \Delta, \bar{\Delta}, \Omega, S, \end{cases} \\
C_1^\phi &= 3Y_\phi^2/5, \quad \text{for each } \phi \text{ with } U(1) \text{ charge } Y_\phi.
\end{aligned}$$

These contributions will eventually translate into soft SUSY breaking terms, and in case of BM the desired effective potential to produce charge-preserving vacuum. However, there is no signal of global parity breakdown and the problem of domain walls persists. In the next section we propose a modification of this standard GMSB to explain parity breaking.

V. CUSTOMIZED GMSB FOR LEFT-RIGHT SYMMETRIC MODELS

The differences required between the soft terms of the left and the right sector for the DW to disappear at a temperature T_D as given in Table I are very small. The reasons for the appearance of this small asymmetry between the left and the right fields are hard to explain since the original theory is parity symmetric. However, we now try to explain the origin of this small difference by focusing on the hidden sector and relating it to SUSY breaking.

For this purpose we assume that the strong dynamics responsible for SUSY breaking also breaks parity, which is then transmitted to the visible sector via the messenger sector and encoded in the soft supersymmetry breaking terms. We implement this idea by introducing two singlet fields X and X' , respectively, even and odd under parity.

$$X \leftrightarrow X, \quad X' \leftrightarrow -X'. \quad (30)$$

The messenger sector superpotential then contains terms

$$\begin{aligned}
W &= \sum_n [\lambda_n X (\Phi_{nL} \bar{\Phi}_{nL} + \Phi_{nR} \bar{\Phi}_{nR}) \\
&\quad + \lambda'_n X' (\Phi_{nL} \bar{\Phi}_{nL} - \Phi_{nR} \bar{\Phi}_{nR})]. \quad (31)
\end{aligned}$$

For simplicity, we consider $n = 1$. The fields $\Phi_L, \bar{\Phi}_L$ and $\Phi_R, \bar{\Phi}_R$ are complete representations of a simple gauge group embedding the L-R symmetry group. Further we require that the fields labeled L get exchanged with fields labeled R under an inner automorphism which exchanges $SU(2)_L$ and $SU(2)_R$ charges, e.g., the charge conjugation operation in $SO(10)$. As a simple possibility we consider the case when $\Phi_L, \bar{\Phi}_L$ (respectively, $\Phi_R, \bar{\Phi}_R$) are neutral under $SU(2)_R$ ($SU(2)_L$). Generalization to other representations is straightforward.

As a result of the dynamical SUSY breaking we expect the fields X and X' to develop nontrivial vev's and F terms and hence give rise to mass scales

$$\Lambda_X = \frac{\langle F_X \rangle}{\langle X \rangle}, \quad \Lambda_{X'} = \frac{\langle F_{X'} \rangle}{\langle X' \rangle}. \quad (32)$$

Both of these are related to the dynamical SUSY breaking scale M_S , however their values are different unless additional reasons of symmetry would force them to be identical. Assuming that they are different but comparable in magnitude we can show that left-right breaking can be achieved simultaneously with SUSY breaking being communicated.

In the proposed model, the messenger fermions receive respective mass contributions

$$m_{f_L} = |\lambda \langle X \rangle + \lambda' \langle X' \rangle|, \quad m_{f_R} = |\lambda \langle X \rangle - \lambda' \langle X' \rangle|, \quad (33)$$

while the messenger scalars develop the masses

$$\begin{aligned}
m_{\phi_L}^2 &= |\lambda \langle X \rangle + \lambda' \langle X' \rangle|^2 \pm |\lambda \langle F_X \rangle + \lambda' \langle F_{X'} \rangle|, \\
m_{\phi_R}^2 &= |\lambda \langle X \rangle - \lambda' \langle X' \rangle|^2 \pm |\lambda \langle F_X \rangle - \lambda' \langle F_{X'} \rangle|. \quad (34)
\end{aligned}$$

We thus have both SUSY and parity breaking communicated through these particles.

As a result the mass contributions to the gauginos of $SU(2)_L$ and $SU(2)_R$ from both the X and X' fields with their corresponding auxiliary parts take the simple form,

$$M_{a_L} = \frac{\alpha_a}{4\pi} \frac{\lambda \langle F_X \rangle + \lambda' \langle F_{X'} \rangle}{\lambda \langle X \rangle + \lambda' \langle X' \rangle} (1 + O(x_L)), \quad (35)$$

where

$$x_L = \frac{\lambda \langle F_X \rangle + \lambda' \langle F_{X'} \rangle}{|\lambda \langle X \rangle + \lambda' \langle X' \rangle|^2} \quad (36)$$

and

$$M_{a_R} = \frac{\alpha_a}{4\pi} \frac{\lambda \langle F_X \rangle - \lambda' \langle F_{X'} \rangle}{\lambda \langle X \rangle - \lambda' \langle X' \rangle} (1 + O(x_R)) \quad (37)$$

with

$$x_R = \frac{\lambda \langle F_X \rangle - \lambda' \langle F_{X'} \rangle}{|\lambda \langle X \rangle - \lambda' \langle X' \rangle|^2}. \quad (38)$$

Here $a = 1, 2, 3$. In turn there is a modification to scalar masses, through two-loop corrections, expressed to leading orders in the x_L or x_R , respectively, by the generic formulas

$$\begin{aligned}
m_{\phi_L}^2 &= 2 \left(\frac{\lambda \langle F_X \rangle + \lambda' \langle F_{X'} \rangle}{\lambda \langle X \rangle + \lambda' \langle X' \rangle} \right)^2 \left[\left(\frac{\alpha_3}{4\pi} \right)^2 C_3^\phi + \left(\frac{\alpha_2}{4\pi} \right)^2 (C_{2L}^\phi) \right. \\
&\quad \left. + \left(\frac{\alpha_1}{4\pi} \right)^2 C_1^\phi \right], \quad (39)
\end{aligned}$$

$$m_{\phi_R}^2 = 2 \left(\frac{\lambda \langle F_X \rangle - \lambda' \langle F_{X'} \rangle}{\lambda \langle X \rangle - \lambda' \langle X' \rangle} \right)^2 \left[\left(\frac{\alpha_3}{4\pi} \right)^2 C_3^\phi + \left(\frac{\alpha_2}{4\pi} \right)^2 (C_{2R}^\phi) + \left(\frac{\alpha_1}{4\pi} \right)^2 C_1^\phi \right], \quad (40)$$

where the values of C_i^ϕ 's are given by Eq. (29). Applying these formulas to Δ and Ω , the parameters m_i^2 's and β_i 's appearing in Eqs. (17) and (19) respectively can be calculated. The differences between the mass squared of the left and right sectors are obtained as

$$\begin{aligned} \delta m_\Delta^2 &= 2 \left[\left(\frac{\lambda \langle F_X \rangle + \lambda' \langle F_{X'} \rangle}{\lambda \langle X \rangle + \lambda' \langle X' \rangle} \right)^2 - \left(\frac{\lambda \langle F_X \rangle - \lambda' \langle F_{X'} \rangle}{\lambda \langle X \rangle - \lambda' \langle X' \rangle} \right)^2 \right] \\ &\quad \times \left\{ \left(\frac{\alpha_2}{4\pi} \right)^2 + \frac{6}{5} \left(\frac{\alpha_1}{4\pi} \right)^2 \right\} \\ &= 2(\Lambda_X)^2 \left[\left(\frac{1 + \tan\gamma}{1 + \tan\sigma} \right)^2 - \left(\frac{1 - \tan\gamma}{1 - \tan\sigma} \right)^2 \right] \\ &\quad \times \left\{ \left(\frac{\alpha_2}{4\pi} \right)^2 + \frac{6}{5} \left(\frac{\alpha_1}{4\pi} \right)^2 \right\} \\ &= 2(\Lambda_X)^2 f(\gamma, \sigma) \left\{ \left(\frac{\alpha_2}{4\pi} \right)^2 + \frac{6}{5} \left(\frac{\alpha_1}{4\pi} \right)^2 \right\}, \end{aligned} \quad (41)$$

where

$$f(\gamma, \sigma) = \left(\frac{1 + \tan\gamma}{1 + \tan\sigma} \right)^2 - \left(\frac{1 - \tan\gamma}{1 - \tan\sigma} \right)^2. \quad (42)$$

We have brought Λ_X out as the representative mass scale and parametrized the ratio of mass scales by introducing

$$\tan\gamma = \frac{\lambda' \langle F_{X'} \rangle}{\lambda \langle F_X \rangle}, \quad \tan\sigma = \frac{\lambda' \langle X' \rangle}{\lambda \langle X \rangle}. \quad (43)$$

Similarly,

$$\delta m_\Omega^2 = 2(\Lambda_X)^2 f(\gamma, \sigma) \left(\frac{\alpha_2}{4\pi} \right)^2. \quad (44)$$

In the models studied here, the ABMRS model will have contribution from both the above kinds of terms. The BM model will have contribution only from the Δ fields.

The contribution to slepton masses is also obtained from Eqs. (39) and (40). This can be used to estimate the magnitude of the overall scale Λ_X to be ≥ 30 TeV [71] from collider limits. Substituting this in the above formulas (41) and (44) we obtain the magnitude of the factor $f(\gamma, \sigma)$ required for cosmology as estimated in Table I. The resulting values of $f(\gamma, \sigma)$ are tabulated in Table II. We see that

TABLE II. Entries in this table are the values of the parameter $f(\gamma, \sigma)$, required to ensure wall disappearance at temperature T_D displayed in the header row. The table should be read in conjunction with Table I, with the rows corresponding to each other.

$T_D/\text{GeV} \sim$	10	10^2	10^3
Adequate ($m^2 - m'^2$)	10^{-7}	10^{-3}	10
Adequate ($\beta_1 - \beta_2$)	10^{-11}	10^{-7}	10^{-3}

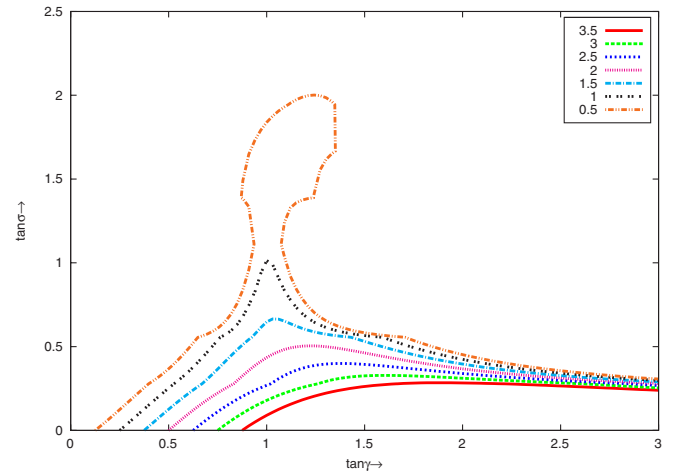


FIG. 1 (color online). The contours represent acceptable values of the parameters for $(m^2 - m'^2) \sim 10^3 \text{ GeV}^2$.

obtaining the values of T_D low compared to the TeV scale requires considerable fine-tuning of f . The natural range of temperature for the disappearance of domain walls therefore remains TeV or higher, i.e., up to a few order of magnitudes lower than the scale at which they form.

Consider for instance $T_D \sim 3 \times 10^2 \text{ GeV}$, which allows $(m^2 - m'^2)$ to range over $\sim 10^2 \text{ GeV}^2$ to 10^3 GeV^2 . Accordingly, we have plotted the range of acceptable values of $\tan\gamma$ and $\tan\sigma$ for two representative values of $(m^2 - m'^2)$ in Figs. 1 and 2. In Fig. 1, we plot the contours of f corresponding to $(m^2 - m'^2) = (2 \pm 1.5) \times 10^3 \text{ GeV}^2$ in steps of 0.5×10^3 . We see that there is a considerable region of parameter space available in this case. In Fig. 2, using $(m^2 - m'^2) \sim 10 \text{ GeV}^2$, we find that

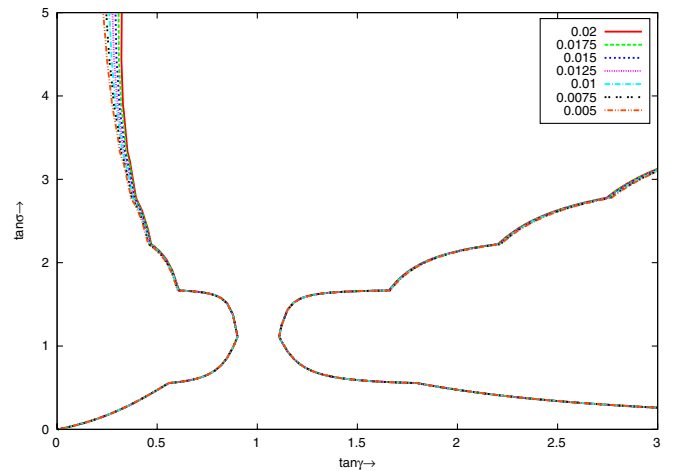


FIG. 2 (color online). The contours represent acceptable values of the parameters for $(m^2 - m'^2) \sim 10 \text{ GeV}^2$. The contours in this case are seen to be overlapping over most of the parameter range. This means extreme sensitivity of one parameter to the value of the other in this range of $(m^2 - m'^2)$ values.

the contours of all the values $(1.25 \pm 0.75) \times 10 \text{ GeV}^2$ collapse to a single curve except in a narrow range of values with $\tan\gamma \sim 0.4$ and $\tan\sigma > 3$. We thus see that in this case, $\tan\gamma$ and $\tan\sigma$ must be highly correlated. The reason for this is that the function $f(\gamma, \sigma)$ is extremely flat except for very restricted parts of the parameter space. This forces an unexpected strong correlation of the scales of F_X , X vevs in the parity conserving sector with $F_{X'}$, X' vevs in the parity violating sectors. While this is specific to the particular scheme we have proposed for the communication of parity violation along with SUSY violation, our scheme we believe is fairly generic and the results may persist for other implementations of this idea.

VI. CONCLUSION

In left-right symmetric models, domain walls generically arise in cosmology. It is necessary to assume the presence of dynamics that eventually signals departure from exact left-right symmetry. In the absence of such a dynamics the Universe would remain trapped in an unacceptable phase. We have explored models where the scale of left-right symmetry and the accompanying gauged $B - L$ symmetry are both low, within a few orders of magnitude of the electroweak scale using 10^4 GeV as a specific example. In earlier work we have obtained bounds on the parity breaking parameters so that domain walls do not conflict with phenomenology, at the same time providing mechanisms for leptogenesis as well as weak inflation [24,25]. The latter is an effective ways of diluting the density of unwanted relics [26]. Wall disappearance is a nonadiabatic phenomenon and could leave behind imprints on primordial gravitational wave background in the range of energy scales we have considered [23].

The possibilities considered in this paper fall into two categories, whether weak inflation is permitted or not. Leptogenesis is permitted in all the cases considered here. Weak inflation becomes possible if the domain walls linger around for a substantial time, dominating the energy density of the Universe for a limited period. The walls are long lived if the pressure difference across the walls is small, as happens if the parity breaking effects are small and the difference in effective potential across the walls is small.

In this paper we have explored the possibility that smallness of the parity breaking effect is related to the indirect supersymmetry breaking effects. To be specific we have studied two viable implementations of left-right symmetry, the ABMRS and BM models discussed in Sec. II, and studied them in the context of gauge mediated supersymmetry breaking as the mechanism. We have explored a variant of the latter mechanism in order to achieve parity breaking to be signaled from within the dynamical supersymmetry breaking sector. Dynamical symmetry breaking effects are associated with strong dynamics and parity is usually susceptible to breaking in their presence.

Our implementation of GMSB contains two singlets X even under parity and X' odd under parity in the hidden sector and coupled to messengers. The scale of left-right symmetry we have explored is 10^4 GeV . We have seen that obtaining the parity breaking effect in this context is natural for values of wall disappearance temperature $T_D \gtrsim 1 \text{ TeV}$. However, T_D lower than the TeV scale generically requires $\frac{\lambda\langle F_{X'} \rangle}{\lambda\langle F_X \rangle}$ and $\frac{\lambda\langle X' \rangle}{\lambda\langle X \rangle}$ to be finely tuned to each other as seen in Table II. We have further explored this effect for two specific ranges of values of $(m^2 - m'^2)$ both of which correspond to $T_D \sim 10^2 \text{ GeV}$. As discussed at the end of Sec. V this confirms the onset of fine-tuning as the value of T_D is lowered. Whether this fine-tuning requirement is generic to other implementations of the main idea of parity breaking communication from the hidden sector remains to be explored. If T_D is as low as 10 GeV or lower, all the cosmological requirements can be met, albeit with extreme fine-tuning.

From Table I we see that the lower bound on T_D required by Ω fields is higher, if the same dynamics determines the soft terms in Δ and Ω effective Lagrangians. Thus Ω vev's determine the T_D . On the other hand from Table II, we see that given a desirable value of T_D the terms in the Δ Lagrangian require less fine-tuning than those in the Ω Lagrangian. Since in the BM model the singlet does not signal any new mass scale, it is the scale of Δ vev's which determines the T_D . For this reason it would be a more natural model from the point of view of cosmology. This result continues to hold if the $W^{(2)}$ terms dropped in Sec. II B are restored.

The ABMRS model also naturally supports long lived domain walls if for some reason the Ω vev scale does not enter the wall disappearance mechanism. One way this could occur is if the SUSY breaking effects were communicated primarily to the Δ sector but not the Ω sector. In a class of models considered in [48,49] SUSY breaking gets communicated by fields charged only under $B - L$ and no other charges. However we have to keep in mind that in this class of models it is difficult to keep the gaugino mass large enough to avoid the existing bounds. Since the Ω are neutral under $B - L$, they would receive the SUSY breaking effects only as higher order effects. In this category of SUSY breaking models ABMRS would require less fine-tuning to ensure a solution of all the cosmological issues studied here.

Our general conclusion is that within the class of models considered, the requirement of resolving the stated issues of cosmology constrains the model greatly. However the natural versions of the model which do not get so constrained are interesting in their own right. While the issues of cosmology may require a separate investigation, this paper has identified a natural implementation of spontaneous parity breaking embedded within gauge mediated supersymmetry breaking.

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APPENDIX: VACUA IN THE BM MODEL

Here we discuss some details of the minimization conditions for the BM model. We begin with the full renormalizable superpotential for the BM model

$$\begin{aligned}
W_{LR} = & \mathbf{h}_l^{(i)} L^T \tau_2 \Phi_i \tau_2 L_c + \mathbf{h}_q^{(i)} Q^T \tau_2 \Phi_i \tau_2 Q_c \\
& + i \mathbf{f}^* L^T \tau_2 \Delta L + i \mathbf{f} L^c \tau_2 \Delta_c L_c + S[\lambda^* \text{Tr} \Delta \bar{\Delta} \\
& + \lambda \text{Tr} \Delta_c \bar{\Delta}_c + \lambda'_{ab} \text{Tr} \Phi_a^T \tau_2 \Phi_b \tau_2 - M_R^2] \\
& + M_\Delta \text{Tr} \Delta \bar{\Delta} + M_\Delta^* \text{Tr} \Delta_c \bar{\Delta}_c \\
& + \mu_{ab} \text{Tr} \Phi_a^T \tau_2 \Phi_b \tau_2 + M_s S^2 + \lambda_s S^3. \quad (A1)
\end{aligned}$$

The resulting expressions for the F terms are

$$F_s = -[\lambda^* \text{Tr} \Delta \bar{\Delta} + \lambda \text{Tr} \Delta_c \bar{\Delta}_c - M_R^2 + 2M_s S + 3\lambda_s S^2]^*, \quad (A2)$$

$$F_\Delta = -2[S\lambda^* \bar{\Delta} + M_\Delta \bar{\Delta}]^*, \quad (A3)$$

$$F_{\bar{\Delta}} = -2[S\lambda^* \Delta + M_\Delta \Delta]^*, \quad (A4)$$

$$F_{\Delta_c} = -2[S\lambda \bar{\Delta}_c + M_\Delta^* \bar{\Delta}_c]^*, \quad (A5)$$

$$F_{\bar{\Delta}_c} = -2[S\lambda \Delta_c + M_\Delta^* \Delta_c]^*. \quad (A6)$$

From these we assemble the potential for the scalar fields in standard notation,

$$\begin{aligned}
V = & |F_s|^2 + |F_\Delta|^2 + |F_{\bar{\Delta}}|^2 + |F_{\Delta_c}|^2 + |F_{\bar{\Delta}_c}|^2 \\
= & |\lambda^* \text{Tr} \Delta \bar{\Delta} + \lambda \text{Tr} \Delta_c \bar{\Delta}_c - M_R^2 + 2M_s S + 3\lambda_s S^2|^2 \\
& + 4|S\lambda^* + M_\Delta|^2 \text{Tr} \bar{\Delta} \bar{\Delta}^\dagger + 4|S\lambda^* + M_\Delta|^2 \text{Tr} \Delta \Delta^\dagger \\
& + 4|S\lambda + M_\Delta^*|^2 \text{Tr} \bar{\Delta}_c \bar{\Delta}_c^\dagger + 4|S\lambda + M_\Delta^*|^2 \text{Tr} \Delta_c \Delta_c^\dagger \\
= & |\lambda^* v_L \bar{v}_L + \lambda v_R \bar{v}_R - M_R^2 + 2M_s S + 3\lambda_s S^2|^2 \\
& + 4|S\lambda^* + M_\Delta|^2 (|\bar{v}_L|^2 + |v_L|^2) \\
& + 4|S\lambda + M_\Delta^*|^2 (|\bar{v}_R|^2 + |v_R|^2). \quad (A7)
\end{aligned}$$

The resulting minimization conditions for the vev's are

$$\begin{aligned}
\frac{\delta V}{\delta S} = & (2M_s + 6\lambda_s S) Q^* + 4\lambda^* (S\lambda^* + M_\Delta)^* \\
& \times (|\bar{v}_L|^2 + |v_L|^2) + 4\lambda (S\lambda + M_\Delta^*)^* \\
& \times (|\bar{v}_R|^2 + |v_R|^2) = 0, \quad (A8)
\end{aligned}$$

$$\frac{\delta V}{\delta v_L} = \lambda^* \bar{v}_L Q^* + 4|S\lambda^* + M_\Delta|^2 v_L^* = 0, \quad (A9)$$

$$\frac{\delta V}{\delta v_L} = \lambda^* v_L Q^* + 4|S\lambda^* + M_\Delta|^2 \bar{v}_L^* = 0, \quad (A10)$$

$$\frac{\delta V}{\delta v_R} = \lambda \bar{v}_R Q^* + 4|S\lambda + M_\Delta^*|^2 v_R^* = 0, \quad (A11)$$

$$\frac{\delta V}{\delta \bar{v}_R} = \lambda v_R Q^* + 4|S\lambda + M_\Delta^*|^2 \bar{v}_R^* = 0, \quad (A12)$$

where

$$Q = \lambda^* v_L \bar{v}_L + \lambda v_R \bar{v}_R - M_R^2 + 2M_s S + \lambda_s S^2. \quad (A13)$$

Thus the desired class of vacua Eq. (11) is obtained provided we ignore the $W^{(2)}$ of Eq. (10) in the text and choose $\langle S \rangle$ to be zero.

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