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Spontaneous parity breaking with broken supersymmetry : cosmological constraint[1]

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Abstract. Unified models incorporating the right handed neutrino in a symmetric way generically possess parity symmetry. If this is broken spontaneously it results in the formation of domain walls in the early Universe, whose persistence is unwanted. A generic mechanism for destabilisation of such walls is a small pressure difference signalled by difference in the free energy across the walls. It is interesting to explore the possibility of such effects in conjunction with the effects that break supersymmetry in a phenomenologically acceptable way. Realising this possibility in the context of several scenarios of supersymmetry breaking results in an upper bound on the scale of spontaneous parity breaking, often much lower than the GUT scale. In the left-right symmetric models studied, the upper bound is no higher than 10^{11}GeV but a scale as low as 10^{5}GeV is acceptable.

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1. LEFT-RIGHT SYMMETRY : A QUICK INTRODUCTION

Chirality seems to be an essential feature of fundamental physics, thereby allowing dynamical generation of fermion masses. However the observed parity violation of the Standard Model (SM) is not warranted by chirality. Discovery of neutrino masses in the past two decades strongly suggests the existence of right handed neutrino states. The resulting parity balanced spectrum of fermions begs a parity symmetric theory and parity violation could then be explained to be of dynamical origin. An interesting fact to emerge is that the see-saw mechanism generically suggests an M_R scale considerably smaller than the scale of coupling constant unification in $SO(10)$. It is therefore appealing to look for left-right symmetry as an intermediate stage in the sequence of symmetry breaking, and explore the possible range of masses acceptable for M_R. The crucial phenomenological question is, could the new symmetries be within the accessible range of the LHC and the colliders of foreseeable future, and hence deserve the name Just Beyond the Standard Model (JBSM)?

Left-right symmetric model[1, 2] needs a Supersymmetric extension as an expedient for avoiding the hierarchy problem. The minimal set of Higgs superfields required, with

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their $SU(3)\otimes SU(2)_L\otimes SU(2)_R\otimes U(1)_{B-L}$ is,

$$
\begin{aligned}
& & \Phi_i = (1,2,2,0), & \qquad i = 1,2, & \\
\Delta = (1,3,1,2), & & \bar{\Delta} = (1,3,1,-2), & \qquad \Delta_c = (1,1,3,-2), & \bar{\Delta}_c = (1,1,3,2), \\
& & \Omega = (1,3,1,0), & \qquad \Omega_c = (1,1,3,0) &
\end{aligned} \tag{1}
$$

and further details of the model can be found in the references. Here we consider a model with the Higgs content which ensures spontaneous parity breaking, preserving electromagnetic charge invariance, and retaining R parity [3]. It contains the two additional triplet Higgs fields introduced in the third line above. We refer to this as ABMRS model. Supersymmetric minima breaking $SU(2)_R$ symmetry are signaled by the ansatz

$$
\langle\Omega_c\rangle = \begin{pmatrix} \omega_c & 0 \\ 0 & -\omega_c \end{pmatrix}, \qquad \langle\Delta_c\rangle = \begin{pmatrix} 0 & 0 \\ d_c & 0 \end{pmatrix}, \tag{2}
$$

In this model, with an enhanced R symmetry, we are lead naturally to a see-saw relation $M_{B-L}^2 = M_{EW}M_R$. This means Leptogenesis is postponed to a lower energy scale closer to M_{EW}. Being generically below 10^9 GeV, this avoids the gravitino mass bound but requires non-thermal leptogensis[4].

For comparison we also take an alternative model to this, considered in [5] where a superfield $S(1,1,1,0)$ also singlet under parity is included in addition to the minimal set of Higgs required. This is referred here as BM model.

2. COSMOLOGY OF BREAKING AND SOFT TERMS

SUSY breaking soft terms emerge below the SUSY breaking scale M_S. We now proceed with the stipulation advanced in [6] that 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 parameterized as follows

$$
\mathscr{L}^1_{soft} = \ m_1^2\mathrm{Tr}(\Delta\Delta^\dagger) + m_2^2\mathrm{Tr}(\bar{\Delta}\bar{\Delta}^\dagger) + \ m_3^2\mathrm{Tr}(\Delta_c\Delta_c^\dagger) + m_4^2\mathrm{Tr}(\bar{\Delta}_c\bar{\Delta}_c^\dagger) \tag{3}
$$

$$
\mathscr{L}^2_{soft} = \alpha_1\mathrm{Tr}(\Delta\Omega\Delta^\dagger) + \alpha_2\mathrm{Tr}(\bar{\Delta}\Omega\bar{\Delta}^\dagger) + \ \alpha_3\mathrm{Tr}(\Delta_c\Omega_c\Delta_c^\dagger) + \alpha_4\mathrm{Tr}(\bar{\Delta}_c\Omega_c\bar{\Delta}_c^\dagger) \tag{4}
$$

$$
\mathscr{L}^3_{soft}\ \beta_1\mathrm{Tr}(\Omega\Omega^\dagger) + \beta_2\mathrm{Tr}(\Omega_c\Omega_c^\dagger) \tag{5}
$$

$$
\mathscr{L}^4_{soft} = S[\gamma_1\mathrm{Tr}(\Delta\Delta^\dagger) + \gamma_2\mathrm{Tr}(\bar{\Delta}\bar{\Delta}^\dagger)] + \ S^*[\gamma_3\mathrm{Tr}(\Delta_c\Delta_c^\dagger) + \gamma_4\mathrm{Tr}(\bar{\Delta}_c\bar{\Delta}_c^\dagger)] \tag{6}
$$

$$
\mathscr{L}^5_{soft} = \ \tilde{\sigma}^2|S|^2 \tag{7}
$$

For ABMRS model the relevant soft terms are given by,

$$
\mathscr{L}_{soft} = \mathscr{L}^1_{soft} + \mathscr{L}^2_{soft} + \mathscr{L}^3_{soft} \tag{8}
$$

TABLE 1. 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/GeV	$\sim$	10	10^2	10^3
$(m^2 - m^{2\prime})/\mathrm{GeV}^2$	$\sim$	10^{-4}	1	10^4
$(\beta_1 - \beta_2)/\mathrm{GeV}^2$	$\sim$	10^{-8}	10^{-4}	1

For BM model the soft terms are given by,

$$\mathscr{L}_{soft} = \mathscr{L}^1_{soft} + \mathscr{L}^4_{soft} + \mathscr{L}^5_{soft} \tag{9}$$

Using the requirement $\delta\rho \sim T_D^4$ we can constrain the differences between the soft terms in the Left and Right sectors [7, 8]. In the BM model the S field does not acquire a vev in the physically relevant vacua and hence the terms in eq.s (6) and (7) do not contribute to the vacuum energy. The terms in eq. (4) are suppressed in magnitude relative to those in eq. (5) 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 $m_3^2 \sim m_4^2 \sim m'^2$ [8] with T_D in the range $10 - 10^3$ GeV [9]. For both the models we have taken the value of the Δ vev's as $d \sim 10^4$ GeV. For ABMRS model additionally we take $\omega \sim 10^6$ GeV. The resulting differences required for successful removal of domain walls are shown in Table 1.

We see from table 1 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 DW as discussed in [10, 8]. After the disappearance of the walls at the scale T_D, electroweak symmetry breaks at a scale $M_{EW} \sim 10^2$ GeV and standard cosmology takes over. In the next section we discuss the implementation of GMSB scenario for these models.

3. TRANSITORY DOMAIN WALLS

Spontaneous parity breaking leads to formation of Domain walls which quickly dominate the energy density of the Universe. It is necessary for recovering standard cosmology that these walls disappear at least before the Big Bang Nucleosynthesis (BBN). In an intrinsically parity symmetric theory difference in the vacua resulting in destabilisation is not permitted. We may seek these effects to have arisen from the hidden sector and communicated along with the messenger fields [11]. Constraints on the hidden sector model and the communication mechanism can be obtained in this way. Here we report on other possibilities.

There are several studies of wall evolution, and an estimate of the temperature at which the walls may destabilise, parametrically expressed in terms of the surface tension of the walls, in turn determined by the parity breaking scale M_R. By equating the terms leading to small symmetry breaking discussed in the previous para with this parametric dependence then gives a bound on M_R.

The dynamics of the walls in a radiation dominated universe is determined by two quantities : [12], *Tension force* $f_T \sim \sigma/R$, where σ is energy per unit area and R is the average scale of radius of curvature, and *Friction force* $f_F \sim \beta T^4$ for walls moving with speed β in a medium of temperature T. The scaling law for the growth of the scale $R(t)$ on which the wall complex is smoothed out, is taken to be $R(t) \approx (G\sigma)^{1/2} t^{3/2}$. Also, $f_F \sim 1/(Gt^2)$ and $f_T \sim (\sigma/(Gt^3))^{1/2}$. Then the pressure difference required to overcome the above forces and destabilise the walls is

$$\delta\rho_{RD} \geq G\sigma^2 \approx \frac{M_R^6}{M_{Pl}^2} \sim M_R^4 \frac{M_R^2}{M_{Pl}^2} \tag{10}$$

The case of matter dominated evolution is relevant to moduli fields copiously produced in generic string inspired models [9] of the Universe. A wall complex formed at temperature $T_i \sim M_R$ is assumed to have first relaxed to being one wall segment per horizon volume. It then becomes comparable in energy density to the ambient matter density, due to the difference in evolution rates, $1/a(t)$ for walls compared to $1/a^3(t)$ for matter. For simplicity also demand that the epoch of equality of the two contributions is the epoch also of instability, so as to avoid dominance by domain walls. Thus we can set $M_{Pl}^{-2} T_D^4 \sim H_{eq}^2 \sim \sigma^{\frac{3}{4}} H_i^{\frac{1}{4}} M_{Pl}^{-3}$. The corresponding temperature permits the estimate of the required pressure difference,

$$\delta\rho_{MD} > M_R^4 \left(\frac{M_R}{M_{Pl}}\right)^{3/2} \tag{11}$$

Thus in this case we find $(M_R/M_{Pl})^{3/2}$ [13], a milder suppression factor than in the radiation dominated case above.

4. PARITY BREAKING FROM PLANCK SUPPRESSED EFFECTS

For a generic neutral scalar field ϕ, the higher dimensional operators that may break parity have the simple form [14] $V_{eff} = \frac{C_5}{M_{Pl}}\phi^5$. But this is only instructional because in realistic theories, the structure and effectiveness of such terms is conditioned by Gauge invariance and supersymmetry and the presence of several scalar species.

One possibility is that the parity breaking operators arise at Planck scale [13]. We shall assume the structure of the symmetry breaking terms as dictated by the Kahler potential formalism and treat the cases of two different kinds of domain wall evolution.

Substituting the VEV's in the effective potential, we get

$$V_{eff}^{R} \sim \frac{a(c_R + d_R)}{M_{Pl}} M_R^4 M_W + \frac{a(a_R + d_R)}{M_{Pl}} M_R^3 M_W^2 \tag{12}$$

and likewise $R \leftrightarrow L$. Hence, with generic coefficients κ, which for naturalness should remain order unity,

$$\delta\rho \sim \kappa^A \frac{M_R^4 M_W}{M_{Pl}} + \kappa'^A \frac{M_R^3 M_W^2}{M_{Pl}} \tag{13}$$

Then equating to $\delta\rho_{RD}$, $\delta\rho_{MD}$ derived above,

$$\kappa_{RD}^{A} > 10^{-10} \left(\frac{M_R}{10^6 \mathrm{GeV}} \right)^2 \tag{14}$$

For M_R scale tuned to 10^9GeV needed to avoid gravitino problem after reheating at the end of inflation, $\kappa_{RD} \sim 10^{-4}$, a reasonable constraint. but requires κ_{RD}^{A} to be $O(1)$ or unnaturally large for the scale of M_R greater than the intermediate scale 10^{11}GeV.

Next,

$$\kappa_{MD}^{A} > 10^{-2} \left(\frac{M_R}{10^6 \mathrm{GeV}} \right)^{3/2}, \tag{15}$$

which seems to be a modest requirement, but taking $M_R \sim 10^9$GeV required to have thermal leptogenesis without the undesirable gravitino production, leads to unnatural $\kappa_{MD} > 10^{5/2}$.

Concluding this section we note that the least restrictive requirement on $\delta\rho$ is $\gtrsim (1\mathrm{MeV})^4$ in order for the walls to not ruin BBN. This requirement gives a *lower* bound on the M_R scale, generically much closer to the TeV scale.

5. 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 are not unnaturally large. However the reasons for appearance of even a small asymmetry between the Left and the Right fields is hard to explain since the original theory is parity symmetric. 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.

$$X \leftrightarrow X, \qquad X' \leftrightarrow -X'. \tag{16}$$

The messenger sector superpotential then contains terms

$$W = \sum_n \left[\lambda_n X \left(\Phi_{nL}\bar{\Phi}_{nL} + \Phi_{nR}\bar{\Phi}_{nR}\right) + \lambda_n' X' \left(\Phi_{nL}\bar{\Phi}_{nL} - \Phi_{nR}\bar{\Phi}_{nR}\right)\right] \tag{17}$$

For simplicity, we consider $n = 1$. The fields Φ_L, $\bar{\Phi}_L$ and Φ_R, $\bar{\Phi}_R$ are complete representations of a simple gauge group embedding the L-R symmetry group. Further we require that the fields labelled L get exchanged with fields labelled 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 Φ_L, $\bar{\Phi}_L$ (respectively, Φ_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

$$\Lambda_X = \frac{\langle F_X \rangle}{\langle X \rangle}, \qquad \Lambda_{X'} = \frac{\langle F_{X'} \rangle}{\langle X' \rangle}. \tag{18}$$

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

$$\begin{aligned} m_{f_L} &= |\lambda \langle X \rangle + \lambda' \langle X' \rangle| \\ m_{f_R} &= |\lambda \langle X \rangle - \lambda' \langle X' \rangle| \end{aligned} \tag{19}$$

while the messenger scalars develop the masses

$$\begin{aligned} m^2_{\phi_L} &= |\lambda \langle X \rangle + \lambda' \langle X' \rangle|^2 \pm |\lambda \langle F_X \rangle + \lambda' \langle F_{X'} \rangle| \\ m^2_{\phi_R} &= |\lambda \langle X \rangle - \lambda' \langle X' \rangle|^2 \pm |\lambda \langle F_X \rangle - \lambda' \langle F_{X'} \rangle| \end{aligned} \tag{20}$$

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,

$$M_{a_L} = \frac{\alpha_a}{4\pi} \frac{\lambda \langle F_X \rangle + \lambda' \langle F_{X'} \rangle}{\lambda \langle X \rangle + \lambda' \langle X' \rangle} \tag{21}$$

and

$$M_{a_R} = \frac{\alpha_a}{4\pi} \frac{\lambda \langle F_X \rangle - \lambda' \langle F_{X'} \rangle}{\lambda \langle X \rangle - \lambda' \langle X' \rangle} \tag{22}$$

upto terms suppressed by $\sim F/X^2$. 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 formulae

$$m_{\phi_L}^2 = 2\left(\frac{\lambda\langle F_X\rangle + \lambda'\langle F_{X'}\rangle}{\lambda\langle X\rangle + \lambda'\langle X'\rangle}\right)^2 \left[\left(\frac{\alpha_3}{4\pi}\right)^2 C_3^\phi + \left(\frac{\alpha_2}{4\pi}\right)^2 (C_{2L}^\phi) + \left(\frac{\alpha_1}{4\pi}\right)^2 C_1^\phi\right] \tag{23}$$

$$m_{\phi_R}^2 = 2\left(\frac{\lambda\langle F_X\rangle - \lambda'\langle F_{X'}\rangle}{\lambda\langle X\rangle - \lambda'\langle X'\rangle}\right)^2 \left[\left(\frac{\alpha_3}{4\pi}\right)^2 C_3^\phi + \left(\frac{\alpha_2}{4\pi}\right)^2 (C_{2R}^\phi) + \left(\frac{\alpha_1}{4\pi}\right)^2 C_1^\phi\right] \tag{24}$$

The resulting difference between the mass squared of the left and right sectors are obtained as

$$\delta m_\Delta^2 = 2(\Lambda_X)^2 f(\gamma,\sigma)\left\{\left(\frac{\alpha_2}{4\pi}\right)^2 + \frac{6}{5}\left(\frac{\alpha_1}{4\pi}\right)^2\right\} \tag{25}$$

where,

$$f(\gamma,\sigma) = \left(\frac{1+\tan\gamma}{1+\tan\sigma}\right)^2 - \left(\frac{1-\tan\gamma}{1-\tan\sigma}\right)^2 \tag{26}$$

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

$$\tan\gamma = \frac{\lambda'\langle F_{X'}\rangle}{\lambda\langle F_X\rangle}, \quad \tan\sigma = \frac{\lambda'\langle X'\rangle}{\lambda\langle X\rangle} \tag{27}$$

Similarly,

$$\delta m_\Omega^2 = 2(\Lambda_X)^2 f(\gamma,\sigma)\left(\frac{\alpha_2}{4\pi}\right)^2 \tag{28}$$

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

The contribution to slepton masses is also obtained from eq.s (23) and (24). This can be used to estimate the magnitude of the overall scale Λ_X to be ≥ 30 TeV [15] from collider limits. Substituting this in the above formulae (25) and (28) we obtain the magnitude of the factor $f(\gamma,\sigma)$ required for cosmology as estimated in table 1. The resulting values of $f(\gamma,\sigma)$ are tabulated in table 2. We see that obtaining the values of T_D low compared to 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., upto a few order of magnitudes lower than the scale at which they form.

Consider for instance $T_D \sim 3\times 10^2$GeV, which allows $(m^2 - m'^2)$ to range over $\sim 10^2$GeV^2 to 10^3 GeV^2. Consider two representative values of $\tan\gamma$ and $\tan\sigma$ for of $(m^2 - m'^2)$. First, $(m^2 - m'^2) = (2\pm 1.5)\times 10^3$GeV^2. This results in sufficient paramour space for the F and X parameters. However when we consider $(m^2 - m'^2) \sim 10$ GeV^2. We find that this requires the two parameters to be fine tuned to each other as $\tan\gamma \sim 0.4$ and $\tan\sigma > 3$. 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.

TABLE 2. 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 conjuction with table 1, with the rows corresponding to each other.

T_D/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}

6. SUPERSYMMETRY BREAKING IN METASTABLE VACUA

The dilemma of phenomenology with broken supersymmetry can be captured in the fate of R symmetry generic to superpotentials [16]. An unbroken R symmetry in the theory is required for SUSY breaking. R symmetry when spontaneously broken leads to R-axions which are unacceptable. If we give up R symmetry, the ground state remains supersymmetric. The solution proposed in [16, 17], is to break R symmetry mildly, governed by a small parameter ε. Supersymmetric vacuum persists, but this can be pushed far away in field space. SUSY breaking local minimum is ensured near the origin, since it persists in the limit $\varepsilon \to 0$. A specific example of this scenario [18] referred to as ISS, envisages $SU(N_c)$ SQCD (UV free) with $N_f(> N_c)$ flavors such that it is dual to a $SU(N_f - N_c)$ gauge theory (IR free) so called magnetic phase, with N_f^2 singlet mesons M and N_f flavors of quarks $q, \tilde{q}$.

Thus we consider a Left-Right symmetric model with ISS mechanism as proposed in [19]. The particle content of the electric theory is $Q_L^a \sim (3,1,2,1,1)$, $\tilde{Q}_L^a \sim (3^*,1,2,1,-1)$ and $Q_R^a \sim (1,3,1,2,-1)$, $\tilde{Q}_R^a \sim (1,3^*,1,2,1)$. where $a = 1, N_f$ with the gauge group G_{33221}. This SQCD has $N_c = 3$, and we need $N_f \geq 4$.

For $N_f = 4$ the dual magnetic theory has Left Right gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ and the effective fields are the squarks and nonet mesons carrying either the $SU(2)_L$ or the $SU(2)_R$ charges. The Left-Right symmetric renormalisable superpotential of this magnetic theory is

$$W_{LR}^0 = h\mathrm{Tr}\phi_L\Phi_L\tilde{\phi}_L - h\mu^2\mathrm{Tr}\Phi_L + h\mathrm{Tr}\phi_R\Phi_R\tilde{\phi}_R - h\mu^2\mathrm{Tr}\Phi_R \tag{29}$$

After integrating out the right handed chiral fields, the superpotential becomes

$$W_L^0 = h\mathrm{Tr}\phi_L\Phi_L\tilde{\phi}_L - h\mu^2\mathrm{Tr}\Phi_L + h^4\Lambda^{-1}\mathrm{det}\Phi_R - h\mu^2\mathrm{Tr}\Phi_R \tag{30}$$

which gives rise to SUSY preserving vacua at

$$\langle h\Phi_R\rangle = \Lambda_m\varepsilon^{2/3} = \mu\frac{1}{\varepsilon^{1/3}} \tag{31}$$

where $\varepsilon = \frac{\mu}{\Lambda_m}$. Thus the right handed sector exists in a metastable SUSY breaking vacuum whereas the left handed sector is in a SUSY preserving vacuum breaking D-parity spontaneously.

We next consider [20] Planck scale suppressed terms that may signal parity breaking

$$W^1_{LR} = f_L \frac{\mathrm{Tr}(\phi_L \Phi_L \tilde{\phi}_L)\mathrm{Tr}\Phi_L}{\Lambda_m} + f_R \frac{\mathrm{Tr}(\phi_R \Phi_R \tilde{\phi}_R)\mathrm{Tr}\Phi_R}{\Lambda_m} + f'_L \frac{(\mathrm{Tr}\Phi_L)^4}{\Lambda_m} + f'_R \frac{(\mathrm{Tr}\Phi_R)^4}{\Lambda_m} \tag{32}$$

The terms of order $\frac{1}{\Lambda_m}$ are given by

$$V^1_R = \frac{h}{\Lambda_m} S_R [f_R(\phi^0_R \tilde{\phi}^0_R)^2 + f'_R \phi^0_R \tilde{\phi}^0_R S^2_R + (\delta^0_R - S_R)^2((\phi^0_R)^2 + (\tilde{\phi}^0_R)^2)] \tag{33}$$

The minimization conditions give $\phi\tilde{\phi} = \mu^2$ and $S^0 = -\delta^0$. Denoting $\langle\phi^0_R\rangle = \langle\tilde{\phi}^0_R\rangle = \mu$ and $\langle\delta^0_R\rangle = -\langle S^0_R\rangle = M_R$, we have

$$V^1_R = \frac{h f_R}{\Lambda_m}(|\mu|^4 M_R + |\mu|^2 M^3_R) \tag{34}$$

where we have also assumed $f'_R \approx f_R$. For $|\mu| < M_R$ Thus the effective energy density difference between the two types of vacua is

$$\delta\rho \sim h(f_R - f_L)\frac{|\mu|^2 M^3_R}{\Lambda_m} \tag{35}$$

Thus for walls disappearing in matter dominated era, we get

$$M_R < |\mu|^{5/9} M^{4/9}_{Pl} \sim 1.3 \times 10^{10}\ \mathrm{GeV} \tag{36}$$

with μ ~Tev. Similarly for the walls disappearing in radiation dominated era,

$$M_R < |\mu|^{10/21} M^{11/21}_{Pl} \sim 10^{11}\ \mathrm{GeV} \tag{37}$$

7. CONCLUSIONS

We have pursued the possibility of left-right symmetric models as Just Beyond Standard Models (JBSM), not possessing a large hierarchy. We also adopt the natural points of view that right handed neutrinos must be included in the JBSM in a symmetric way and that the required parity breaking to match low energy physics arises from spontaneous breakdown. The latter scenario is often eschewed due to the domain walls it entails in the early Universe. We turn the question around to ask given that the domain walls occur, what physics could be responsible for their successful removal without jeopardising naturalness.

We do not advance any preferred way to provide the small asymmetry required to get rid of the domain walls. However it is interesting to correlate the possibility that these small effects may be correlated to the supersymmetry breaking. We have considered three models along these lines. One in which the hidden sector breaking of supersymmetry is at a low energy, and mediated by a gauge sector. Another in which the generic scale of supersymmetry breaking is at Planck scale and the breaking effects are conveyed purely through Planck scale suppressed terms. Finally we have also considered a possible implementation of the scenarios in which the supersymmetry breaking is not in a hidden sector but occurs due to a metastable vacuum protected from decay by a large suppression of tunnelling.

The general message seems to be that the parity breaking scale in any case is not warranted to be as high as required for a full unification in $SO(10)$ and further, several scenarios suggest that left-right symmetry as the larger package incorporating the SM may be within the reach of future colliders.

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