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Why PeV scale left–right symmetry is a good thing

URJIT A YAJNIK

Physics Department, Indian Institute of Technology Bombay, Mumbai 400 076, India
E-mail: yajnik@iitb.ac.in

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Abstract. Left–right symmetric gauge theory presents a minimal paradigm to accommodate massive neutrinos with all the known conserved symmetries duly gauged. The work presented here is based on the argument that the see-saw mechanism does not force the new right-handed symmetry scale to be very high, and as such some of the species from the spectrum of the new gauge and Higgs bosons can have masses within a few orders of magnitude of the TeV scale. The scale of the left–right parity breaking in turn can be sequestered from the Planck scale by supersymmetry. We have studied several formulations of such just beyond Standard Model (JBSM) theories for their consistency with cosmology. Specifically, the need to eliminate phenomenologically undesirable domain walls gives many useful clues. The possibility that the exact left–right symmetry breaks in conjunction with supersymmetry has been explored in the context of gauge mediation, placing restrictions on the available parameter space. Finally, we have also studied a left–right symmetric model in the context of metastable supersymmetric vacua and obtained constraints on the mass scale of right-handed symmetry. In all the cases studied, the mass scale of the right-handed neutrino M_R remains bounded from above, and in some of the cases the scale 10^9 GeV favourable for supersymmetric thermal leptogenesis is disallowed. On the other hand, PeV scale remains a viable option, and the results warrant a more detailed study of such models for their observability in collider and astroparticle experiments.

Keywords. Left–right symmetry; neutrino mass; domain walls; supersymmetry; metastable vacua.

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1. Introduction: The case for low-scale right-handed symmetry

Chirality seems to be an essential feature of fundamental physics, thereby allowing dynamical generation of fermion masses. However, this did not require the world to be parity asymmetric, the way it is manifested in the Standard Model (SM). Indeed, the discovery of neutrino masses[1–5] 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. Let us take stock of what principles we could put to use in a 'down upwards' guesswork in energy scales.

- Gauge principle provides massless force carriers
- Chirality provides massless matter
- Scalars can signal spontaneous symmetry breakdown
  - provide masses through a universal mechanism
  - provide naturally substantial amount of CP violation.
- Supersymmetry provides the most comprehensive and elegant way of sequestering mass scales signalled by scalars.

Any extension we seek could first be guided by these principles. To be specific, we assume the restriction that any exact internal symmetry should be gauged, and any scalar and fermionic fields introduced must be charged under at least one of these gauge symmetries. The left–right class of models are interesting from this point of view and are explored as the first rung of a 'down-upwards' unification. The two other alternatives not pursued by us are: (1) $SO(10)$ unification with grand desert, as for instance proposed in [6–8] and (2) the SM Higgs boson so instrumental in providing a perturbative description of weak interactions is secretly a member of a strongly coupled theory [9–13].

As baryogenesis began to be ruled out in SM [14,15] and leptogenesis [16] became severely constrained in thermal $SO(10)$ scenarios [17,18], our early observation [19] was that the see-saw mechanism [20–24] generically easily permits an M_R scale as low as 10^6 GeV, considerably smaller than the scales of coupling constant unification. Secondly, that low scale non-thermal leptogenesis is consistent with the wash-out constraints from low-scale right-handed neutrinos. It is therefore appealing to look for left–right symmetry [25] as an intermediate stage in the sequence of symmetry breaking, and explore the possible range of masses $M_{\rm R}$ acceptable for the heavy right-handed neutrinos. In the hope that the new symmetries are within the reach of the LHC and several proposed future colliders, we may call it just beyond the Standard Model (JBSM).

A number of recent works have focussed on this question and studied it in conjunction with correlated signatures such as with baryogenesis or leptogenesis, and also additionally dark matter production [26–33] and gravitational waves [34]. There are also planned or ongoing experiments such as non-accelerator exploration of N–$\bar{N}$ oscillations [35], role of right-handed currents in the electric dipole moments of light nuclei [36] and hadronic CP violation [37], correlated signatures in neutrinoless double beta decay [38,39] on the one hand and collider signatures [31,40,41] on the other hand. Secondly, a number of viable scenarios for low-scale leptogenesis exist, such as resonant leptogenesis [42,43] as also the generically supersymmetric Affleck–Dine scenario [44,45]. We considered specific models [46] (non-supersymmetric) and [47] (supersymmetric version) of non-thermal leptogenesis, but it is presumably possible for the other scenarios such as referred in this paragraph also to be implementable within the parameter space of the left–right models to be considered here. Assuming that this is possible within the PeV scale, we have studied the consistency of some of the intrinsic features of the proposed left–right symmetric models with cosmology.

2. Left–right symmetry: A supersymmetric revival

Left–right symmetric model [22,48] needs a supersymmetric extension as an expedient for avoiding the hierarchy problem. The minimal set of Higgs superfields required, with their $SU(3)\otimes SU(2)_{\rm L}\otimes SU(2)_{\rm R}\otimes U(1)_{\rm B-L}$ is,

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

and further details of the model can be found in the references.

There is an awkward impasse with this model, namely we would like to retain supersymmetry down to the TeV scale. So the first stage of gauge symmetry breaking has to respect supersymmetry. If we choose the parameters of the superpotential to ensure spontaneous parity breaking, then either the electromagnetic gauge invariance or the R parity has to be sacrificed [49,50]. The first of these is unacceptable, and the second entails a requirement of inelegant fixes. This problem was elegantly resolved in [51] with further developments in [52,53]. It contains the two additional triplet Higgs fields

$$\Omega = (1,3,1,0), \qquad \Omega_c = (1,1,3,0). \tag{2}$$

We refer to this briefly as ABMRS model. Supersymmetric minima breaking $SU(2)_{\rm R}$ symmetry are signalled by the ansatz

$$\langle\Omega_c\rangle = \begin{pmatrix} \omega_c & 0 \\ 0 & -\omega_c \end{pmatrix}, \qquad \langle\Delta_c\rangle = \begin{pmatrix} 0 & 0 \\ d_c & 0 \end{pmatrix}. \tag{3}$$

In this model, with an enhanced R symmetry, we are led naturally to a see-saw relation $M_{\rm B-L}^2 = M_{\rm EW}M_{\rm R}$. This means leptogenesis is postponed to a lower energy scale closer to $M_{\rm EW}$. Being generically below 10^9 GeV, this avoids the gravitino mass bound but requires non-thermal leptogensis [54].

For comparison, we also take an alternative model to this, considered in [55] where a superfield $S(1,1,1,0)$ also singlet under parity is included in addition to the minimal set of Higgs of eq (1). This is referred to here as the BM model.

3. Transitory domain walls

A generic problem long recognized with the model with exact left–right symmetry is that of cosmological domain walls. The effective potential of the theory is unable to give preference to whether $SU(2)_{\rm L}$ is broken first or $SU(2)_{\rm R}$. There is therefore a competition to the SM vacuum from a 'parity-flipped SM' vacuum, with the low-energy effective gauge group $SU(3)_{\rm c}\otimes SU(2)_{\rm R}\otimes U(1)_{\tilde{Y}}$ where $\tilde{Y} = T_{\rm L}^3 + \frac{1}{2}(B-L)$. Spontaneous parity breaking leads to the 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. In an intrinsically parity-symmetric theory, difference in the vacua resulting in destabilization is not permitted. In the following we first obtain the quantitative requirement on the possible asymmetry so that standard cosmology is ensured. In the next two subsections, we propose mechanisms that could be the source of such asymmetry without causing serious damage to the basic assumption of exact left–right symmetry.

There are several studies of wall evolution, and an estimate of the temperature at which the walls may destabilize, 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 paragraph 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 [56]: Tension force $f_T \sim \sigma/R$, where σ is the 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 destabilize the walls is

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

The case of matter-dominated evolution is relevant to moduli fields copiously produced in generic string-inspired models [57] 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. T_{stab} is taken to be the temperature when the walls begin to be unstable and less dominant. Thus, we can set $M_{Pl}^{-2} T_{destab}^4 \sim H_{eq}^2 \sim \sigma^{3/4} H_i^{1/4} M_{Pl}^{-3}$. The corresponding temperature permits the estimate of the required pressure difference,

$$\delta\rho_{MD} > M_R^4 \left(\frac{M_R}{M_{Pl}}\right)^{3/2}. \tag{5}$$

Thus, in this case we find $(M_R/M_{Pl})^{3/2}$ [58], as a milder suppression factor than in the radiation-dominated case above.

3.1 *Soft terms as a source of lifting parity degeneracy*

One source that could provide the required pressure $\delta\rho$ can be sought in a generic neutral scalar field ϕ. It provides the higher dimensional operator that may break parity and which has the simple form [59] $V_{eff} = (C_5/M_{Pl})\phi^5$. The cosmological requirement then constrains the coefficient C_5. In realistic theories, there are several scalar fields entering such terms, and the structure of the latter is conditioned by gauge invariance and supersymmetry. Here we study a possible source of such asymmetric terms without sacrificing the symmetries of the superpotential, viz., in the soft supersymmetry breaking terms. The soft terms which arise in the two models, ABMRS and BM respectively, may be parametrized as follows:

$$\mathcal{L}_{soft}^1 = m_1^2 \mathrm{Tr}(\Delta\Delta^\dagger) + m_2^2 \mathrm{Tr}(\bar{\Delta}\bar{\Delta}^\dagger) + m_3^2 \mathrm{Tr}(\Delta_c \Delta_c^\dagger) + m_4^2 \mathrm{Tr}(\bar{\Delta}_c \bar{\Delta}_c^\dagger) \tag{6}$$

$$\mathcal{L}_{soft}^2 = \alpha_1 \mathrm{Tr}(\Delta\Omega\Delta^\dagger) + \alpha_2 \mathrm{Tr}(\bar{\Delta}\Omega\bar{\Delta}^\dagger) + \alpha_3 \mathrm{Tr}(\Delta_c\Omega_c\Delta_c^\dagger) + \alpha_4 \mathrm{Tr}(\bar{\Delta}_c\Omega_c\bar{\Delta}_c^\dagger) \tag{7}$$

$$\mathcal{L}_{soft}^3 = \beta_1 \mathrm{Tr}(\Omega\Omega^\dagger) + \beta_2 \mathrm{Tr}(\Omega_c\Omega_c^\dagger) \tag{8}$$

$$\mathcal{L}_{soft}^4 = S[\gamma_1 \mathrm{Tr}(\Delta\Delta^\dagger) + \gamma_2 \mathrm{Tr}(\bar{\Delta}\bar{\Delta}^\dagger)] + S^*[\gamma_3 \mathrm{Tr}(\Delta_c\Delta_c^\dagger) + \gamma_4 \mathrm{Tr}(\bar{\Delta}_c\bar{\Delta}_c^\dagger)] \tag{9}$$

$$\mathcal{L}_{soft}^5 = \tilde{\sigma}^2 |S|^2. \tag{10}$$

For ABMRS model, the relevant soft terms are given by

$$\mathcal{L}_{soft} = \mathcal{L}_{soft}^1 + \mathcal{L}_{soft}^2 + \mathcal{L}_{soft}^3. \tag{11}$$

For BM model the soft terms are given by

$$\mathcal{L}_{soft} = \mathcal{L}_{soft}^1 + \mathcal{L}_{soft}^4 + \mathcal{L}_{soft}^5. \tag{12}$$

It can be shown [60] that the adequate requirement to evade the domain wall problem is to demand $\delta\rho \sim T_{destab}^4$, where $\delta\rho$ is the difference in the effective potentials across the domain walls and T_{destab} is the destabilization temperature, at which this difference becomes the dominant force acting on the walls. At a later epoch, the walls decay at temperature T_D. If the disappearance of the walls is prompt, the two temperatures would be comparable. Even when the epochs are separated in time, the entropy generated due to final wall disappearance may raise the temperature, bringing it closer to T_{destab}. In the absence of a detailed model, $T_D \sim T_{destab}$. To obtain the observed Universe we need that T_D remains higher than the big-bang nucleosynthesis (BBN) scale of a few MeV. This is the requirement imposed to constrain the differences between the soft terms in the left and right sectors [47,61]. In the BM model, the S field does not acquire a vacuum expectation value (vev) in the physically relevant vacua and hence the terms in eqs (9) and (10) do not contribute to the vacuum energy. The terms in eq. (7) are suppressed in magnitude relative to those in eq. (8) 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 $m_3^2 \sim m_4^2 \sim m'^2$ [47] with T_D in the range 10–10^3 GeV

Table 1. Differences in values of soft supersymmetry breaking parameters for a range of domain wall disappearance temperature values T_D. The differences signify the extent of parity breaking.

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

[57]. 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 the 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 cosmology. This is because the lower bound on the wall decay 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.

Between the epoch of destabilization of the DWs and their disappearance, leptogenesis occurs due to preferential motion of the DW to select SM as the low energy theory, as discussed in [46,47]. After the disappearance, i.e., complete decay of the walls, also assumed to occur at the scale T_D, electroweak symmetry breaks at a scale $M_{EW} \sim 10^2$ GeV and standard cosmology takes over.

Since the soft terms are associated with supersymmetry breaking, they may be assumed to arise from the same mechanism that breaks supersymmetry. In §4 we discuss the implementation of gauge-mediated supersymmetry breaking (GMSB) scenario for these models and treat the soft terms to have arisen from the hidden sector and communicated along with the messenger fields [62]. Constraints on the hidden sector model and the communication mechanism can be obtained in this way.

3.2 *Parity breaking from Planck-suppressed effects*

The soft terms studied above are tied to the scale of supersymmetry breaking. Another mechanism to look into, without incurring violence to the symmetries of the superpotential, is to assume that the parity-breaking operators arise at the Planck scale [58]. Writing the lowest order terms in the superpotential suppressed by the Planck scale, and using the Kähler potential formalism, we obtain the expectation value of the effective potential of the scalar fields after substituting the vacuum expectation values of the relevant fields as follows, whose details can be found in [58].

$$V_{\text{eff}}^{\text{R}} \sim \frac{a(c_{\text{R}} + d_{\text{R}})}{M_{\text{Pl}}} M_{\text{R}}^4 M_{\text{W}} + \frac{a(a_{\text{R}} + d_{\text{R}})}{M_{\text{Pl}}} M_{\text{R}}^3 M_{\text{W}}^2 \tag{13}$$

and likewise $R \leftrightarrow L$. Here the constants a, c_R etc. are dimensionless. Hence, by regrouping the above coefficients into coefficients κ etc., which for naturalness should remain order unity,

$$\delta\rho \sim \kappa^A \frac{M_{\text{R}}^4 M_{\text{W}}}{M_{\text{Pl}}} + \kappa'^A \frac{M_{\text{R}}^3 M_{\text{W}}^2}{M_{\text{Pl}}}. \tag{14}$$

It is interesting to study this relation with κ coefficients $O(1)$ and $\delta\rho \sim (10\ \text{MeV})^4$, compatible with BBN. It then leads without any further assumptions to a lower bound on M_R to be the same as SM scale, leaving all higher scales open to being physically viable.

Now inserting these answers in the expectations derived from cosmological dynamics of DW, viz., $\delta\rho_{\text{RD}}$, $\delta\rho_{\text{MD}}$ derived in eqs (4) and (5), we obtain

$$\kappa_{\text{RD}}^A > 10^{-10} \left(\frac{M_{\text{R}}}{10^6\ \text{GeV}} \right)^2. \tag{15}$$

Thus, only minuscule values of κ coefficients suffice to allow PeV scale M_R. At the higher end, $M_R \lesssim 10^9$ GeV needed to avoid gravitino problem after reheating at the end of inflation remains viable with $\kappa_{RD} \sim 10^{-4}$, and also M_R of the order of the intermediate scale 10^{11} GeV remains marginally viable with $\kappa_{RD}^A \sim O(1)$. But scales higher than that are not tolerated by naturalness.

Next, if we consider wall disappearance during moduli-dominated regime, we find

$$\kappa_{\text{MD}}^A > 10^{-2} \left(\frac{M_{\text{R}}}{10^6\ \text{GeV}} \right)^{3/2}, \tag{16}$$

which can be seen to be a rather strict requirement barely allowing M_R higher than our preferred PeV scale. In particular, taking $M_R \sim 10^9$ GeV required to have thermal leptogenesis without the undesirable gravitino production, leads to unnatural values of $\kappa_{MD} > 10^{5/2}$.

4. 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 because the theory as adopted in §2 is parity symmetric.

We now try to explain the origin of this small difference by focussing on the hidden sector, and relating it to supersymmetry (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 is 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.

$$X \leftrightarrow X, \qquad X' \leftrightarrow -X'. \tag{17}$$

The messenger sector superpotential then contains terms

$$W = \sum_n \left[\lambda_n X\left(\Phi_{n\mathrm{L}}\bar{\Phi}_{n\mathrm{L}} + \Phi_{n\mathrm{R}}\bar{\Phi}_{n\mathrm{R}}\right) + \lambda'_n X'\left(\Phi_{n\mathrm{L}}\bar{\Phi}_{n\mathrm{L}} - \Phi_{n\mathrm{R}}\bar{\Phi}_{n\mathrm{R}}\right)\right]. \tag{18}$$

For simplicity, we consider $n = 1$. The fields Φ_{L}, $\bar{\Phi}_{\mathrm{L}}$ and Φ_{R}, $\bar{\Phi}_{\mathrm{R}}$ 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 $SU(2)_{\mathrm{L}}$ and $SU(2)_{\mathrm{R}}$ charges, e.g. the charge conjugation operation in $SO(10)$. As a simple possibility, we consider the case when Φ_{L}, $\bar{\Phi}_{\mathrm{L}}$ (respectively, Φ_{R}, $\bar{\Phi}_{\mathrm{R}}$) are neutral under $SU(2)_{\mathrm{R}}$ ($SU(2)_{\mathrm{L}}$). Generalization to other representations is straightforward.

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

$$\Lambda_X = \frac{\langle F_X\rangle}{\langle X\rangle}, \qquad \Lambda_{X'} = \frac{\langle F_{X'}\rangle}{\langle X'\rangle}. \tag{19}$$

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

$$m_{f_{\mathrm{L}}} = |\lambda\langle X\rangle + \lambda'\langle X'\rangle|$$
$$m_{f_{\mathrm{R}}} = |\lambda\langle X\rangle - \lambda'\langle X'\rangle| \tag{20}$$

while the messenger scalars develop the masses

$$m^2_{\phi_{\mathrm{L}}} = |\lambda\langle X\rangle + \lambda'\langle X'\rangle|^2 \pm |\lambda\langle F_X\rangle + \lambda'\langle F_{X'}\rangle|$$
$$m^2_{\phi_{\mathrm{R}}} = |\lambda\langle X\rangle - \lambda'\langle X'\rangle|^2 \pm |\lambda\langle F_X\rangle - \lambda'\langle F_{X'}\rangle|. \tag{21}$$

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

As a result, the mass contributions to the gauginos of $SU(2)_{\mathrm{L}}$ and $SU(2)_{\mathrm{R}}$ from both the X and X' fields with their corresponding auxiliary parts take the simple form

$$M_{a_{\mathrm{L}}} = \frac{\alpha_a}{4\pi}\frac{\lambda\langle F_X\rangle + \lambda'\langle F_{X'}\rangle}{\lambda\langle X\rangle + \lambda'\langle X'\rangle} \tag{22}$$

and

$$M_{a_{\mathrm{R}}} = \frac{\alpha_a}{4\pi}\frac{\lambda\langle F_X\rangle - \lambda'\langle F_{X'}\rangle}{\lambda\langle X\rangle - \lambda'\langle X'\rangle} \tag{23}$$

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 x_{L} or x_{R} respectively, by the generic formulae

$$m^2_{\phi_{\mathrm{L}}} = 2\left(\frac{\lambda\langle F_X\rangle + \lambda'\langle F_{X'}\rangle}{\lambda\langle X\rangle + \lambda'\langle X'\rangle}\right)^2 \times\left[\left(\frac{\alpha_3}{4\pi}\right)^2 C_3^{\phi} + \left(\frac{\alpha_2}{4\pi}\right)^2 (C_{2\mathrm{L}}^{\phi}) + \left(\frac{\alpha_1}{4\pi}\right)^2 C_1^{\phi}\right] \tag{24}$$

$$m^2_{\phi_{\mathrm{R}}} = 2\left(\frac{\lambda\langle F_X\rangle - \lambda'\langle F_{X'}\rangle}{\lambda\langle X\rangle - \lambda'\langle X'\rangle}\right)^2 \times\left[\left(\frac{\alpha_3}{4\pi}\right)^2 C_3^{\phi} + \left(\frac{\alpha_2}{4\pi}\right)^2 (C_{2\mathrm{R}}^{\phi}) + \left(\frac{\alpha_1}{4\pi}\right)^2 C_1^{\phi}\right]. \tag{25}$$

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

$$\delta m^2_{\Delta} = 2(\Lambda_X)^2 f(\gamma,\sigma)\left\{\left(\frac{\alpha_2}{4\pi}\right)^2 + \frac{6}{5}\left(\frac{\alpha_1}{4\pi}\right)^2\right\}, \tag{26}$$

where

$$f(\gamma,\sigma) = \left(\frac{1+\tan\gamma}{1+\tan\sigma}\right)^2 - \left(\frac{1-\tan\gamma}{1-\tan\sigma}\right)^2. \tag{27}$$

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

$$\tan\gamma = \frac{\lambda'\langle F_{X'}\rangle}{\lambda\langle F_X\rangle}, \quad \tan\sigma = \frac{\lambda'\langle X'\rangle}{\lambda\langle X\rangle}. \tag{28}$$

Similarly,

$$\delta m^2_{\Omega} = 2(\Lambda_X)^2 f(\gamma,\sigma)\left(\frac{\alpha_2}{4\pi}\right)^2. \tag{29}$$

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 (24) and (25). This can be used to estimate the magnitude of the overall scale Λ_X to be $\geq$30 TeV based on [63] from LEP limits, which may not have

Table 2. Entries in this table are the values of the parameter $f(\gamma, \sigma)$, required to ensure wall decay at temperature T_D displayed in the header row. The table should be read in conjunction with table 1, with the rows corresponding to each other.

T_D (GeV)	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}

changed significantly even in the light of LHC data [64, 65]. Substituting this in the above formulae (26) and (29) 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 the natural range of temperature for the disappearance of domain walls therefore remains TeV or higher, i.e., upto a few orders of magnitudes lower than the scale at which they form.

5. Supersymmetry breaking in metastable vacua

The dilemma of phenomenology with broken supersymmetry can be captured in the fate of R symmetry generic to superpotentials [66]. 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 [66], 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 $\epsilon \to 0$. A specific example of this scenario [67–69] referred to as ISS, envisages $SU(N_c)$ SQCD with $N_f > N_c$ flavours of quarks $q, \tilde{q}$ which is UV-free, such that it is dual to a $SU(N_f - N_c)$ gauge theory which is IR-free, the so-called magnetic phase, with N_f^2 single mesons M.

Thus, we consider a left–right symmetric model with ISS mechanism as proposed in [70]. They proposed the electric gauge theory to be based on the gauge group $SU(3)_L \times SU(3)_R \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ (in short G_{33221}) where $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ is the gauge group of the usual left–right models and $SU(3)_{L,R}$ is the new strongly coupled gauge sector introduced. The dual description similar to the original ISS model gives rise to $SU(2)_R$ broken metastable vacua inducing spontaneous SUSY breaking simultaneously.

The particle content of the electric theory is

$$Q_L^a \sim (3, 1, 2, 1, 1), \quad \tilde{Q}_L^a \sim (3^*, 1, 2, 1, -1),$$
$$Q_R^a \sim (1, 3, 1, 2, -1), \quad \tilde{Q}_R^a \sim (1, 3^*, 1, 2, 1),$$

where $a = 1, N_f$ and the numbers in brackets correspond to the transformations of the fields under the gauge group G_{33221}. This model has $N_c = 3$ and hence to have a Seiberg dual [67] magnetic theory, number of flavours should be $N_f \geq 4$. For $N_f = 4$ the dual magnetic theory will have the gauge symmetry of the usual left–right models $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ and the following particle content

$$\phi_L^a(2, 1, -1), \quad \tilde{\phi}_L^a(2, 1, 1)$$
$$\phi_R^a(1, 2, 1), \quad \tilde{\phi}_R^a(1, 2, -1)$$
$$\Phi_L \equiv \mathbf{1} + \mathrm{Adj}_L = \begin{pmatrix} \frac{1}{\sqrt{2}}(S_L + \delta_L^0) & \delta_L^+ \\ \delta_L^- & \frac{1}{\sqrt{2}}(S_L - \delta_L^0) \end{pmatrix}$$
$$\Phi_R \equiv \mathbf{1} + \mathrm{Adj}_R = \begin{pmatrix} \frac{1}{\sqrt{2}}(S_R + \delta_R^0) & \delta_R^+ \\ \delta_R^- & \frac{1}{\sqrt{2}}(S_R - \delta_R^0) \end{pmatrix}. \tag{30}$$

The left–right symmetric renormalizable superpotential of this magnetic theory is

$$W_{LR}^0 = h\mathrm{Tr}\phi_L\Phi_L\tilde{\phi}_L - h\mu^2\mathrm{Tr}\Phi_L + h\mathrm{Tr}\phi_R\Phi_R\tilde{\phi}_R - h\mu^2\mathrm{Tr}\Phi_R. \tag{31}$$

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

$$W_L^0 = h\mathrm{Tr}\phi_L\Phi_L\tilde{\phi}_L - h\mu^2\mathrm{Tr}\Phi_L + h^4\Lambda^{-1}\mathrm{det}\Phi_R - h\mu^2\mathrm{Tr}\Phi_R \tag{32}$$

which gives rise to SUSY preserving vacua at

$$\langle h\Phi_R\rangle = \Lambda_m\epsilon^{2/3} = \mu\frac{1}{\epsilon^{1/3}}, \tag{33}$$

where $\epsilon = \mu/\Lambda_m$. 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 [71] Planck scale suppressed terms that may signal parity breaking

$$W_{LR}^1 = f_L\frac{\mathrm{Tr}(\phi_L\Phi_L\tilde{\phi}_L)\mathrm{Tr}\Phi_L}{\Lambda_m} + f_R\frac{\mathrm{Tr}(\phi_R\Phi_R\tilde{\phi}_R)\mathrm{Tr}\Phi_R}{\Lambda_m} + f_L'\frac{(\mathrm{Tr}\Phi_L)^4}{\Lambda_m} + f_R'\frac{(\mathrm{Tr}\Phi_R)^4}{\Lambda_m}. \tag{34}$$

The terms of order $1/\Lambda_m$ are given by

$$V_R^1 = \frac{h}{\Lambda_m}S_R[f_R(\phi_R^0\tilde{\phi}_R^0)^2 + f_R'\phi_R^0\tilde{\phi}_R^0S_R^2 + (\delta_R^0 - S_R)^2((\phi_R^0)^2 + (\tilde{\phi}_R^0)^2)]. \tag{35}$$

The minimization conditions give $\phi\tilde{\phi} = \mu^2$ and $S^0 = -\delta^0$. Denoting $\langle\phi_{\rm R}^0\rangle = \langle\tilde{\phi}_{\rm R}^0\rangle = \mu$ and $\langle\delta_{\rm R}^0\rangle = -\langle S_{\rm R}^0\rangle = M_{\rm R}$, we have

$$V_{\rm R}^1 = \frac{hf_{\rm R}}{\Lambda_m}(|\mu|^4 M_{\rm R} + |\mu|^2 M_{\rm R}^3), \tag{36}$$

where we have also assumed $f'_{\rm R} \approx f_{\rm R}$. For $|\mu| < M_{\rm R}$, thus, the effective energy density difference between the two types of vacua is

$$\delta\rho \sim h(f_{\rm R} - f_{\rm L})\frac{|\mu|^2 M_{\rm R}^3}{\Lambda_m}. \tag{37}$$

Now we set $\mu \sim$ TeV, the scale of supersymmetry breaking and as QCD cannot be expected to break parity, assume that the parity breaking terms are essentially of Planck scale origin. Using this self-consistent requirement, for walls disappearing in matter-dominated era, we get

$$M_{\rm R} < |\mu|^{5/9} M_{\rm Pl}^{4/9} \sim 1.3 \times 10^{10}\ {\rm GeV} \tag{38}$$

with $\mu \sim$ TeV. Similarly, for the walls disappearing in radiation-dominated era,

$$M_{\rm R} < |\mu|^{10/21} M_{\rm Pl}^{11/21} \sim 10^{11}\ {\rm GeV}. \tag{39}$$

Thus, in both cases we get an upper bound on $M_{\rm R}$ which can at best be an intermediate scale, while the PeV scale remains eminently viable.

6. 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 with local gauge symmetric interactions 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 jeopardizing naturalness.

Seeking origins of parity breaking that are consistent with the principles outlined in the Introduction, we proceed to correlate this breaking to supersymmetry breaking. We have considered three models along these lines. One involves gauge-mediated supersymmetry breaking, another assumes Planck scale breaking and the third relies on metastable vacua for supersymmetry breaking. The operators permissible in these scenarios are then constrained by the cosmological requirements on the dynamics of domain wall disappearance. In all the cases studied, the mass scale of the right-handed neutrino $M_{\rm R}$ remains bounded from above, and in some of the cases the scale 10^9 GeV favourable for supersymmetric thermal leptogenesis is disallowed. On the other hand, PeV scale remains a viable option.

The possibility that such a low-energy model may be embedded successfully in the semisimple $SO(10)$ at a high scale has been explored separately [72,73]. The general message seems to be that the parity breaking scale is not warranted to be as high as the grand unification scale and further, several scenarios suggest that left–right symmetry as the just beyond Standard Model package incorporating the SM may be within the reach of future colliders.

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