Cosmology in a supersymmetric model with gauged $B - L$

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Abstract

We consider salient cosmological features of a supersymmetric model which is left-right symmetric and therefore possessing gauged $B - L$ symmetry. The requirement of breaking parity and also obtaining charge preserving vacua introduces some unique features to this model (MSLRM), resulting in a preference for non-thermal leptogenesis. Assuming that the model preserves TeV scale supersymmetry, we show that the vacuum structure generically possesses domain walls, which can serve two important purposes. They can signal a secondary inflation required to remove unwanted relics such as gravitino and moduli and also generate lepton asymmetry by a mechanism similar to electroweak baryogenesis. The requirement of disappearance of domain walls imposes constraints on the soft parameters of the theory, testable at the TeV scale. We also propose an alternative model with spontaneous parity violation (MSLRP). Incorporating the same cosmological considerations in this case entails constraints on a different set of soft parameters.

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I. INTRODUCTION

The minimal natural requirement that each of the three right handed neutrino states $\nu_{lR}$ be a doublet partner of the corresponding right handed charged lepton $l_R$ results in a universal left-right symmetric gauge theory by demand, but also leads to gauging of an important exact global symmetry of the standard model, viz., $B - L$. It also results in a satisfactory embedding for the electroweak hypercharge. A natural explanation of the very small observed neutrino masses then resides in the seesaw mechanism, with prediction of heavy Majorana neutrino states, $N_i, i = 1, 2, 3$, whose masses remain model dependent but substantially higher than the electroweak scale. While elegant, seesaw mechanism predicts a new high scale which gives rise to a hierarchy, further complicating the Higgs sector whose standard model manifestation is also poorly understood. Inclusion of supersymmetry (SUSY) however improves the situation, stabilizing hierarchies of mass scales that lie above the SUSY breaking scale. We assume the most optimistic value for the SUSY breaking scale, being the TeV scale without disturbing the standard model. In this paper we study what has been called the minimal supersymmetric left-right symmetric model (MSLRM) with the gauge group $SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$. We explore the possibility for the left-right symmetric scale to be low, a few orders of magnitude removed from the TeV scale. At a higher energy scale the model may turn out to be embedded in the supersymmetric $SO(10)$.

The most minimal supersymmetric left-right symmetric model considered [1, 2] fails to provide spontaneous breakdown of parity. The addition of a parity odd singlet was considered in [3], however this makes the charge preserving vacuum energetically disfavored [1]. Aulakh et al. [4, 5] discussed the inclusion of two new Higgs triplet fields $\Omega$ and $\Omega_c$. In this model, $SU(2)_R$ first breaks to its subgroup $U(1)_R$, at a scale $M_R$, without affecting the $U(1)_{B-L}$. At a lower scale $M_{B-L}$, $SU(3)_c \otimes SU(2)_L \otimes U(1)_R \otimes U(1)_{B-L}$ breaks to $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$. In this scheme parity is spontaneously broken while preserving electromagnetic charge invariance. However, due to parity invariance of the original theory, the phenomenologically unacceptable phase $SU(3)_c \otimes SU(2)_R \otimes U(1)_Y$ is energetically degenerate with $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$. An important consequence of this [6] which we pursue here, is that in the early Universe, domain walls (DW) form at the scale $M_R$.

Domain walls arising due to topological reasons [7] play a crucial role in early cosmology.
If stable, they invalidate the model. Independently, the over-abundance of the gravitino and moduli fields which are typically regenerated after the primordial inflation [8] is a generic problem in most SUSY models [8, 9]. These two problematic ingredients of early cosmology have a happy bearing on each other. It has been shown [6, 10, 11] that if DW dominate the evolution of the universe for a limited duration, the associated secondary inflation removes the gravitino and other dangerous relic fields like moduli. For this scenario to work, it is crucial that the DW are metastable, with a decay temperature $T_D$ which must be larger than $\sim 10$ MeV in order to not interfere with big bang nucleosynthesis (BBN). While a scenario employing transient DW seems to lack the possibility of direct verification, it has been recently pointed out [12] that the upcoming space based gravitational wave detectors may be able to detect the stochastic background arising at such phase transitions. In this paper we shall examine two models, MSLRM and MSLRP, for the possibility of these requirements to be satisfied. The main result is constraints on soft SUSY breaking parameters in the Higgs potentials.

Another issue of cosmology is that this class of models does not favor thermal leptogenesis for an intriguing reason. $B - L$ asymmetry in the form of fermion chemical potential is guaranteed to remain zero in the model until the gauged $B - L$ symmetry breaks spontaneously. As we shall see, a generic consequence of the model is a relation among the various mass scales $M^2_{B-L} \simeq M_W M_R$, where $M_W$ is the electroweak scale. Thermal leptogenesis requires $M_{B-L}$ to be larger than $10^{11}$-10$^{13}$ GeV, which pushes $M_R$ into the Planck scale in light of the above formula. A more optimistic constraint $M_{B-L} > 10^9$GeV [13, 14] requires left-right symmetry to be essentially grand unified theory. More realistic scenarios therefore demand leptogenesis to be non-thermal in this class of models, either through bubble walls of a first order phase transition at the electroweak scale, or DW of the parity breaking phase transition. It has been shown [15, 16] 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 non-thermal. This places the modest bound $M_1 > 10^4$GeV, on the mass of the lightest of the heavy Majorana neutrinos.

The paper is organized as follows. We first review the MSLRM in sec. II. We then outline the main features of cosmology in such a model in sec. III. In sec. IV we introduce the new model MSLRP and discuss the differences to cosmology that can arise. In sec. V we identify the condition for gravitino dilution and its consistency with the $M_R$ scale acceptable
in these models. In secs. VI and VII we write down the permissible soft terms in the two models and obtain the constraints on the parameters determining the safe disappearance of domain walls. Sec. VIII contains summary of conclusions and an outlook.

II. DEGENERATE VACUA OF MSLRM

The minimal supersymmetric left-right symmetric model (MSLRM) [4] contains quark and lepton superfields, one set for each generation, with their quantum numbers under $SU(3)_c, SU(2)_L, SU(2)_R, U(1)_{B-L}$ respectively given by

$$Q = (3, 2, 1, 1/3), \quad Q_c = (3^*, 1, 2, -1/3),
L = (1, 2, 1, -1), \quad L_c = (1, 1, 2, 1).$$

where we have suppressed the generation index. The minimal set of Higgs superfields required is,

$$\Phi_i = (1, 2, 2, 0), \quad i = 1, 2,$$
$$\Delta = (1, 3, 1, 2), \quad \bar{\Delta} = (1, 3, 1, -2),$$
$$\Delta_c = (1, 1, 3, -2), \quad \bar{\Delta}_c = (1, 1, 3, 2),$$

where the bidoublet is doubled so that the model has non-vanishing Cabibo-Kobayashi-Maskawa matrix. The number of triplets is doubled to have anomaly cancellation. Under discrete parity symmetry the fields are prescribed to transform as,

$$Q \leftrightarrow Q^*_c, \quad L \leftrightarrow L^*_c, \quad \Phi_i \leftrightarrow \Phi_i^\dagger,$$
$$\Delta \leftrightarrow \Delta^*_c, \quad \bar{\Delta} \leftrightarrow \bar{\Delta}^*_c.$$

It has been demonstrated in [1] that this minimal generalization to supersymmetry results in a vacuum with $\langle \Delta \rangle = \langle \bar{\Delta} \rangle = \langle \Delta_c \rangle = \langle \bar{\Delta}_c \rangle$. Thus the electroweak scale physics would no longer be chiral. This was also demonstrated to be a general feature for a class of related models. It was proposed to cure this problem by introducing a parity odd singlet [3]. However this results in electromagnetic charge violating vacua. This conclusion in turn can be avoided, but at the cost of violating $R$ parity.

These problems were circumvented in [4], i.e., spontaneous parity breaking, preserving electromagnetic charge invariance, and retaining $R$ parity, can all be achieved by introducing
two new triplet Higgs fields with the following charges.

\[ \Omega = (1, 3, 1, 0), \quad \Omega_c = (1, 1, 3, 0). \]  \hfill (4)

Under parity \( \Omega \leftrightarrow \Omega_c^* \). We shall assume that supersymmetry is broken only at the electroweak scale. Thus at higher scales we look for supersymmetry preserving vacua. Such vacua of the theory are obtained by imposing F-flatness and D-flatness conditions. The conditions can be found in [4] and are similar to the case we have worked out in the Appendix A for the modified version of this model we propose later in sec. IV. The conditions lead to the following set of vacuum expectation value (vev’s) for the Higgs fields as one of the possibilities,

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

At the scale \( M_R \), when \( \Omega, \Omega_c \) acquire vev, \( SU(2)_R \) is broken to \( U(1)_R \). At a lower scale \( M_{B-L} \), the vev of the triplet Higgs fields breaks \( SU(3)_c \otimes SU(2)_L \otimes U(1)_R \otimes U(1)_{B-L} \) to \( SU(3)_c \otimes SU(2)_L \otimes U(1)_Y \). Thus, at TeV scale, the model breaks exactly to the minimal supersymmetric standard model (MSSM). It was shown [4] that this scheme of breaking preserved electromagnetic charge invariance and parity was spontaneously broken. Further, although the \( \Delta \) fields signal \( B - L \) breaking, \( R \) parity is preserved due to the fact that

\[ R = (-1)^{3(B-L)+2s} \]  \hfill (6)

and the \( \Delta \) field vev’s violate \( B - L \) by at least 2 units.

We now observe that the \( D \) and \( F \) flatness conditions imposed above also permit a vacuum preserving the \( SU(2)_R \otimes U(1)_L \otimes U(1)_{B-L} \) symmetry which is energetically degenerate with the one identified in eqns. (5). The alternative set of vev’s is given by

\[ \langle \Omega_c \rangle = 0, \quad \langle \Delta_c \rangle = 0, \quad \langle \bar{\Delta}_c \rangle = 0, \quad \langle \Omega \rangle = \begin{pmatrix} \omega & 0 \\ 0 & -\omega \end{pmatrix}, \quad \langle \Delta \rangle = \begin{pmatrix} 0 & 0 \\ d & 0 \end{pmatrix}, \quad \langle \bar{\Delta} \rangle = \begin{pmatrix} 0 & \bar{d} \\ 0 & 0 \end{pmatrix}. \]  \hfill (7)

This is only to be expected from the \( L \leftrightarrow R \) symmetry of the model. As we discuss below this leads to the formation of domain walls and we must have a mechanism in the model for the removal of such walls.
III. PATTERN OF SYMMETRY BREAKING, COSMOLOGY AND BARYOGENESIS

The solutions for the vev’s from the conditions (5) are

\[ \omega = -\frac{m_\Delta}{a} \equiv -M_R, \]
\[ d = \bar{d} = \left( \frac{2m_\Delta m_\Omega}{a^2} \right)^{1/2} \equiv M_{B-L} \]

As discussed earlier the requirement of obtaining parity breakdown is \( M_R \gg M_{B-L} \), which can be arranged by choosing \( m_\Delta \gg m_\Omega \) which leads to the relation \( M_{B-L}^2 \simeq M_R M_\Omega \). We also adopt the proposal of [5] and assume that \( m_\Omega \simeq m_W \). This is natural, provided the \( m_\Omega \) originates from the soft SUSY breaking sector, which also means in turn that \( m_\Omega \) is comparable to gravitino mass \( m_{3/2} \). Thus we effectively have the relation

\[ M_{B-L}^2 \simeq M_R M_W. \] 

The most optimistic value for \( M_{B-L} \) is \( \sim 10^4 \) GeV with corresponding \( M_R \sim 10^6 \) GeV. On the other hand, the largest value of \( M_R \) is the intermediate scale value \( \sim \sqrt{M_P M_W} \sim 10^{10} \) GeV, beyond which non-renormalizable Planck scale corrections begin to be significant. This corresponds to \( M_{B-L} \sim 10^6 \) GeV. We shall thus be interested in this range of values for \( M_R \) and \( M_{B-L} \), however the lower values make the model amenable to investigation at colliders.

The model contains DW separating the phases identified in (5) and (7). The DW must be unstable for cosmological reasons. The required asymmetry between the two types of vacua has to arise dynamically. Since this is not admissible in the superpotential, it must arise from the soft terms [6]. This means that the mechanism inducing the soft terms must cause a bias between the two types of vacua. In a gauge mediated supersymmetry breaking (GMSB) scenario, the hidden sector or the messenger sector or both must cause a distinction between the two kinds of vacua. On the other hand in a gravity mediated scenario it is permissible to violate the discrete symmetry due to gravitational effects.

An alternative possibility for parity breaking is discussed in section IV. In both cases, below the TeV scale the theory is effectively MSSM. Due to the coupling of the bidoublets \( \Phi_i \) to \( \Omega_c \), one pair of doublets becomes heavy and only one pair of \( SU(2)_L \) doublets remains light.
The discussion so far shows that the model has a characteristic cosmological history, with the $SU(2)_R$ breaking first when the $\Omega$ fields acquire a vev. This is accompanied by the formation of domain walls, with SUSY still preserved. The DW come to dominate the energy density of the Universe and cause the onset of secondary inflation. As the temperature reduces, at the scale $M_{B-L}$, the triplets $\Delta$ acquire vev and the heavy neutrinos obtain Majorana mass. In principle the DW persist down to electroweak scale and lower, possibly upto a decay temperature $T_D$ in the range 10 MeV-10 GeV [11]. The secondary inflation helps to remove unwanted relics viz., gravitinos and moduli fields regenerated after the GUT or Planck scale inflation.

The final disappearance of DW completes the phase transition which commenced at $M_R$. We designate this temperature as $T_D$. The decay process of DW gives rise to entropy production and reheating in a model dependent way. In several models the energy scale of the first order phase transition can be such as to leave behind stochastic gravitational wave background detectable at upcoming space based gravitational wave experiments as pointed out by [12].

Since both $B$ and $L$ are effectively conserved below the electroweak scale, baryogenesis requires that the reheat temperature should be significantly higher than that of the electroweak phase transition. The reheat temperature after secondary inflation ($T_R^s$) however, should not be so high as to regenerate the unwanted relics. Thus, in this model it is required that the secondary inflation has $10^9 > T_R^s > 10^2$ GeV. The upper limit on $T_R^s$ is seen to be easily satisfied for most of the range of values for $M_R \sim 10^6 - 10^{10}$ GeV in this class of models as seen from the energy scales of symmetry breaking.

Furthermore, we saw that the natural scale of $M_{B-L}$ is $10^4 - 10^6$ GeV. Thus thermal leptogenesis is also disfavored and must proceed via one of several possible non-thermal mechanisms. Low scale leptogenesis with special attention to left-right symmetry and supersymmetric mechanisms has already received attention in several works [17, 18, 19, 20, 21, 22, 23].

Specifically the conditions outlined in [17] are easily seen to be satisfied in the present class of models. (i) The phase transition at the $M_R$ scale is necessarily first order due to formation of DW in turn inevitable due to parity invariance of the underlying theory. This ensures the out-of-equilibrium condition necessary for leptogenesis. (ii) Since the DW decay must return the universe to $SU(2)_L$ preserving rather than $SU(2)_R$ preserving phase their is a directionality to DW motion resulting in time reversal violation. (iii) The presence of
several complex scalar fields in the model allows the formation of CP violating condensate in the core of the DW. (iv) This CP violating phase can enter the Dirac mass matrix of the neutrinos streaming through the DW giving rise to leptogenesis governed by the “classical force” mechanism [17, 24, 25, 26, 27]. (v) The lepton asymmetry gets partially converted to baryon asymmetry due to the action of the sphalerons.

IV. SUPERSYMMETRIC LEFT-RIGHT MODEL WITH SPONTANEOUS PARITY BREAKING (MSLR/ P)

We have argued above the need for GMSB or gravity induced soft terms in MSLRM that can lift the symmetry between left and right vacua. While it is easy to see that only one of the two phases can survive, understanding of the selection of the observed vacuum requires additional details in the Planck scale model.

A more appealing option has been considered by Chang et al. [28] where spontaneous breaking of parity is implemented within the Higgs structure of the theory. The model is based on $SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L} \otimes P$, where $P$ denotes a parity symmetry. A gauge singlet field $\eta$ is introduced which is odd under $P$, viz., $\eta \leftrightarrow -\eta$. The potential of the Higgs fields contains a term

$$V_{\eta\Delta} \sim M_{\eta}(\Delta_L^\dagger \Delta_L - \Delta_R^\dagger \Delta_R)$$

(10)

so that if $\eta$ acquires a vev at a high scale $M_P \gg M_W$, $SU(2)_R$ is not broken, but the effective masses of the $\Delta_L$ and $\Delta_R$ become different and the $L \leftrightarrow R$ symmetry appears explicitly broken.

A direct implementation of this idea in supersymmetric theory however, would lead us back to the model of Kuchimanchi and Mohapatra [1] and charge breaking vacua. We propose an alternative model based on the group $SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L} \otimes P$, with a pair of triplets as in [4]. However, unlike [4] the Higgs triplets $\Omega, \Omega_c$ are odd under the parity symmetry. Specifically,

$$Q \leftrightarrow Q_c, \quad L \leftrightarrow L_c, \quad \Phi_i \leftrightarrow \Phi_i^T,$$

$$\Delta \leftrightarrow \Delta_c, \quad \bar{\Delta} \leftrightarrow \bar{\Delta}_c, \quad \Omega \leftrightarrow -\Omega_c.$$

(11)

We dub this model MSLRP. The superpotential consistent with this parity is given by the
The following expression, with a few essential differences from the superpotential of [4].

\[
W_{LR} = h_l^{(i)} L^T \tau_2 \Phi_i \tau_2 L_c + h_q^{(i)} Q^T \tau_2 \Phi_i \tau_2 Q_c + i f L^T \tau_2 \Delta L + i f L^c T \tau_2 \Delta_c L_c \\
+ m_\Delta \text{Tr} \Delta \bar{\Delta} + m_\Delta \text{Tr} \Delta_c \bar{\Delta}_c + \frac{m_\Omega}{2} \text{Tr} \Omega^2 + \frac{m_\Omega}{2} \text{Tr} \Omega_c^2 \\
+ \mu_{ij} \text{Tr} \Phi_i \tau_2 \Phi^T_j \tau_2 - \alpha_{ij} \text{Tr} \Omega_c \Phi_i \tau_2 \Phi^T_j \tau_2 ,
\]

(12)

where color and flavor indices have been suppressed. Further, \(h_q^{(i)} = h_q^{(i)\dagger}\), \(h_l^{(i)} = h_l^{(i)\dagger}\), \(\mu_{ij} = \mu_{ji} = \mu_{ij}^*\), \(\alpha_{ij} = -\alpha_{ji}\). Finally, \(f, h\) are real symmetric matrices with respect to flavor indices.

The \(F\) and \(D\) flatness conditions derived from this superpotential are presented in appendix A. However, the effective potential for the scalar fields which is determined from modulus square of the \(D\) terms remains the same as for the MSLRM at least for the form of the ansatz of the vev’s we have chosen. As such the resulting solution for the vev’s remains identical to eq. (8). The difference in the effective potential shows up in the soft terms as will be shown later. Due to soft terms, below the scale \(M_R\) the effective mass contributions to \(\Delta_c\) and \(\bar{\Delta}_c\) become larger than those of \(\Delta\) and \(\bar{\Delta}\). The cosmological consequence of this is manifested after the \(M_{B-L}\) phase transition when the \(\Delta\)’s become massive. Unlike MSLRM where the DW are destabilized only after the soft terms become significant, i.e., at the electroweak scale, the DW in this case become unstable immediately after \(M_{B-L}\). Leptogenesis therefore commences immediately below this scale and the scenario becomes qualitatively different from that for the MSLRM. This is a subject for future study. In the present work we focus on the removal of unwanted relics and safe exit from DW dominated secondary inflation.

V. DILUTION OF GRAVITINOS

It is reasonable to assume that any primordial abundance of gravitinos has been diluted by the primordial inflation. The gravitinos with potential consequences to observable cosmology are generated entirely after reheating of the universe \((T_R \sim 10^9\text{GeV})\) subsequent to primordial inflation. Detailed calculations [8] show that the gravitino number density \((n_{3/2})\) at a low temp \(T_f\) is given by

\[
n_{3/2}(T_f) = 3.35 \times 10^{-12} \ T_9^{\text{max}} \ T_f^3 \times (1 - 0.018 \ln T_9^{\text{max}}) ,
\]

(13)
where, $T_{9}^{\text{max}} = T/10^9$ GeV

Here we constrain the dynamics governing the DW by requiring that they dilute this regenerated gravitino abundance adequately. The possible values of $M_R$ and $M_{B-L}$ are discussed below eq. (9). Here onwards we shall assume $M_R = 10^6$ GeV and $M_{B-L} = 10^4$ GeV as an example. Putting $T_f \approx M_R = 10^6$ GeV in the above eqn. we get,

$$n_{3/2}(T_f) = 3.35 \times 10^6 (\text{GeV})^3 \equiv n_{3/2}^b ,$$

where, $n_{3/2}^b$ is the gravitino number density at the beginning of secondary inflation at the temperature $T_f = 10^6$ GeV. The number density of gravitino ($n_{3/2}$), during this time, decreases as $R^{-3}$, where $R$ is the scale factor. Therefore we have,

$$R_e = R_b \left( \frac{n_{3/2}^b}{n_{3/2}^e} \right)^{1/3},$$

where, $R_b$ ($R_e$) and $n_{3/2}^b$ ($n_{3/2}^e$) are the scale factor and gravitino density at beginning (end) of secondary inflation.

The best constraint that can be imposed on the gravitinos, produced after primordial inflation, comes from the fact that entropy produced due to decay of gravitino shouldn’t disturb the delicate balance of light nuclei abundance [8, 29]. This constraint is given by,

$$\frac{m_{3/2}\beta n_{3/2}}{n_e E_*} \lesssim 1.$$  

Here, $f$ is the fraction of the gravitino number density ($n_{3/2}$) that dumps its entropy in the universe, and is taken to be $f = 0.8$ [8], $m_{3/2} = 100$ GeV is the mass of the gravitino and $T$ is the scale of BBN taken to be $\sim 1$ MeV. Finally, $E_* = 100$ MeV [29], and $\beta$ is a mildly temperature dependent parameter, with numerically determined value 1.6, which lead to the estimate

$$n_{3/2} \lesssim 1.66 \delta_B \times 10^{-13} (\text{GeV})^3 \equiv n_{3/2}^e ,$$

where $\delta_B [\equiv n_B/n_\gamma]$ is the baryon to photon ratio and $n_\gamma = (2\zeta(3)/\pi^2) T^3$ is the photon number density. Using the Wilkinson Microwave Anisotropic Probe (WMAP) data ($\delta_B = (6.1^{+0.3}_{-0.2}) \times 10^{-10}$) [30] We find that

$$n_{3/2}^e = 1.013 \times 10^{-22} (\text{GeV})^3.$$  

So from eq (15) we have,

$$R_e = R_b \times 3.2 \times 10^9.$$  

Therefore, number of $e$ foldings required is

$$N_e = \ln \left( \frac{R_e}{R_b} \right) = \ln 10^9 \simeq 20 .$$  \hspace{1cm} (20)

This agrees with the observation by [10] that a secondary inflation can dilute the moduli and gravitinos sufficiently that no problem results for cosmology. As pointed out by [9] there can also be more than one secondary inflation to effectively reduce the moduli/gravitino number density. Here we shall assume this to the only secondary inflation sufficient for diluting the gravitino density.

Finally, a handle on the explicit symmetry breaking parameters of the two models can be obtained by noting that there should exist sufficient wall tension for the walls to disappear before a desirable temperature scale $T_D$. It has been observed in [31] that the free energy density difference $\delta \rho$ between the vacua, which determines the pressure difference across a domain wall should be of the order

$$\delta \rho \sim T_D^4$$  \hspace{1cm} (21)

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

In the scenarios we consider in the next two sections, the DW form at the higher scale $M_R$. However at this point the thermal vacua on the two sides of the walls are both equal in free energy. At a lower scale, due to additional symmetry breaking taking effect, more field condensates become a part of the DW structure. After such changes the free energy balance across the walls can change and a net $\delta \rho$ can arise. We shall now see specifically the sources of such changes and relate the resulting $\delta \rho$ to the parameters in appropriate potentials.

VI. REMOVAL OF DOMAIN WALLS : MSLRM

The possible source for breaking the parity symmetry of the MSLRM lies in soft terms with the assumption that the hidden sector, or in case of GMSB also perhaps the messenger sector does not obey the parity of the visible sector model. For gravity mediated breaking this can be achieved in a natural way since a discrete symmetry can be generically broken by gravity effects. We present the possible soft terms for MSLRM below.

$$\mathcal{L}_{soft} = \alpha_1 \text{Tr}(\Omega \Omega^\dagger) + \alpha_2 \text{Tr}(\bar{\Omega} \bar{\Omega}^\dagger) + \alpha_3 \text{Tr}(\Delta_c \Omega_c \Delta_c^\dagger) + \alpha_4 \text{Tr}(\bar{\Delta}_c \Omega_c \bar{\Delta}_c^\dagger) \hspace{1cm} (22)$$

$$+ m_1^2 \text{Tr}(\Delta \Delta^\dagger) + m_2^2 \text{Tr}(\bar{\Delta} \bar{\Delta}^\dagger) + m_3^2 \text{Tr}(\Delta_c \Delta_c^\dagger) + m_4^2 \text{Tr}(\bar{\Delta}_c \bar{\Delta}_c^\dagger) \hspace{1cm} (23)$$

$$+ \beta_1 \text{Tr}(\Omega \Omega^\dagger) + \beta_2 \text{Tr}(\bar{\Omega} \bar{\Omega}^\dagger) .$$  \hspace{1cm} (24)
\[ T_D = 10 \text{ GeV} \quad T_D = 10^2 \text{ GeV} \quad T_D = 10^3 \text{ GeV} \]

| \( m^2 - m^2' \) | \( 10^{-4} \text{ GeV}^2 \) | \( 1 \text{ GeV}^2 \) | \( 10^4 \text{ GeV}^2 \) |
| \( \beta_1 - \beta_2 \) | \( 10^{-8} \text{ GeV}^2 \) | \( 10^{-4} \text{ GeV}^2 \) | \( 1 \text{ GeV}^2 \) |

**TABLE I**: Asymmetry in parameters, for a range of \( T_D \), signifying magnitude of explicit parity breaking

The contributions to the free energy difference \( \delta \rho \) i.e. difference between the left and right sector, can now be estimated from the above Lagrangian. It is natural to consider \( \alpha_1 \sim \alpha_2 \) and \( \alpha_3 \sim \alpha_4 \). In this case it can be shown that the use of eq. (21) does not place a severe constraint on the \( \alpha_i \)'s. For the rest of the soft terms [(23) and (24)] we have respectively, in obvious notation

\[ \delta \rho_{\Delta} = [m_1^2 \text{Tr}(\Delta \Delta^\dagger) + m_2^2 \text{Tr}(\bar{\Delta} \bar{\Delta}^\dagger)] - [m_3^2 \text{Tr}(\Delta_c \Delta_c^\dagger) + m_4^2 \text{Tr}(\bar{\Delta}_c \bar{\Delta}_c^\dagger)] = 2(m^2 - m^2')d^2, \quad (25) \]

\[ \delta \rho_{\Omega} = \beta_1 \text{Tr}(\Omega \Omega^\dagger) - \beta_2 \text{Tr}(\Omega_c \Omega_c^\dagger) = 2(\beta_1 - \beta_2)\omega^2, \quad (26) \]

where we have considered \( m_1^2 \lesssim m_2^2 \equiv m^2, m_3^2 \lesssim m_4^2 \equiv m'^2 \). The vev’s of neutral component of \( \Delta(\Delta_c) \) and \( \Omega(\Omega_c) \) are \( d(d_c) \) and \( \omega(\omega_c) \). Here we have assumed that \( d_c \sim d \) and \( \omega_c \sim \omega \).

Using the constraint (21) in the eqns. (25), (26), we can determine the differences between the relevant soft parameters for a range of permissible values of \( T_D \). In Table I we have taken \( d \sim 10^4 \) GeV, \( \omega \sim 10^6 \) GeV and \( T_D \) in the range 100 MeV – 10 GeV [11]. The above differences between the values in the left and right sectors is a lower bound on the soft parameters and is very small. Larger values would be acceptable to low energy phenomenology. However if we wish to retain the connection to the hidden sector, and have the advantage of secondary inflation we would want the differences to be close to this bound. As pointed out in [31, 32] an asymmetry \( \sim 10^{-12} \) is sufficient to lift the degeneracy between the two sectors.
VII. REMOVAL OF DOMAIN WALLS : MSLRP

In this model parity breaking is achieved spontaneously within the observable sector below the scale $M_R$ at which the $\Omega$ fields acquire vev’s. However the breaking is not manifested in the vacuum till the scale $M_{B-L}$ where the $\Delta$ fields acquire vev’s. For simplicity we assume that the hidden sector responsible for SUSY breaking does not contribute parity breaking terms. This is reasonable since even if the hidden sector breaks this parity the corresponding effects are suppressed by the higher scale of breaking and in the visible sector the parity breaking effects are dominated by the explicit mechanism proposed. Thus at a scale above $M_R$ but at which SUSY is broken in the hidden sector we get induced soft terms respecting this parity. Accordingly, for the Higgs sector the parameters can be chosen such that

$$L_{soft} = \alpha_1 \text{Tr}(\Delta \Omega \Delta^\dagger) - \alpha_2 \text{Tr}(\bar{\Delta} \Omega \bar{\Delta}^\dagger) - \alpha_1 \text{Tr}(\Delta_c \Omega_c \Delta^\dagger_c) + \alpha_2 \text{Tr}(\bar{\Delta}_c \Omega_c \bar{\Delta}^\dagger_c)$$

$$+ m_1^2 \text{Tr}(\Delta \Delta^\dagger) + m_2^2 \text{Tr}(\bar{\Delta} \bar{\Delta}^\dagger) + m_1^2 \text{Tr}(\Delta_c \Delta^\dagger_c) + m_2^2 \text{Tr}(\bar{\Delta}_c \bar{\Delta}^\dagger_c)$$

$$+ \beta \text{Tr}(\Omega \Omega^\dagger) + \beta \text{Tr}(\Omega_c \Omega_c^\dagger) . \quad (27)$$

These terms remain unimportant at first due to the key assumption leading to MSSM as the effective low energy theory. The SUSY breaking effects become significant only at the electroweak scale. However, below the scale $M_R$, $\Omega$ and $\Omega_c$ acquire vev’s given by eq. (5) or (7). Further, below the scale $M_{B-L}$ the $\Delta$ fields acquire vev’s and become massive. The combined contribution from the superpotential and the soft terms to the $\Delta$ masses now explicitly encodes the parity breaking,

$$\mu_\Delta^2 = M_\Delta^2 + \alpha_1 \omega, \quad \mu_\Delta_c^2 = M_\Delta^2 - \alpha_1 \omega,$$

$$\mu_\Delta^2 = M_\Delta^2 + \alpha_2 \omega, \quad \mu_\Delta_c^2 = M_\Delta^2 - \alpha_2 \omega. \quad (28)$$

where $M_\Delta^2$ is the common contribution from the superpotential. The difference in free energy across the domain wall is now dominated by the differential contribution to the $\Delta$ masses

$$\delta \rho_\alpha \equiv 2(\alpha_1 + \alpha_2) \omega d^2; \quad (29)$$

where we have considered $\omega_c \sim \omega$, $d \sim \bar{d} \sim d_c \sim \bar{d}_c$. Now using eq (21) for a range of temperatures ($T_D \sim 10^2$ GeV $- 10^4$ GeV), determines the corresponding range of values of coupling constants as

$$(\alpha_1 + \alpha_2) \sim 10^{-6} - 10^2 \text{ GeV}, \quad (30)$$

where we have considered $|\omega| \simeq M_R$, $|d| \simeq M_{B-L}$. 

VIII. CONCLUSIONS

Supersymmetric left-right model is an appealing model to consider beyond the MSSM due to its natural inclusion of right handed neutrino and gauged $B - L$ symmetry. There is a generic problem with building this kind of models due to their inability to preserve electromagnetic charge invariance together with preserving $R$ parity. MSLRM solves the problem by breaking $SU(2)_R$ at a higher scale and $U(1)_{B-L}$ at a lower scale. A generic consequence of this improvement is a relation $M_{B-L}^2 \simeq M_R M_W$. If $M_R$ is at most of intermediate scale value $10^{10}\text{GeV}$ we get the natural range of values $10^4\text{GeV} - 10^6\text{GeV}$ for the $M_{B-L}$. A generic problem of this class of models remains the need to select the $SU(2)_L$ as the low energy gauge group as against $SU(2)_R$. We propose a new model MSLR/P with spontaneous parity breaking arising in Higgs sector. Both models face the same problems as MSSM with regard to baryogenesis and there is a strong case for leptogenesis to be non-thermal due to the low natural scale for $M_{B-L}$. The pattern of breaking and associated cosmological events in the two classes of models are summarized in table II.

Parity invariance of the underlying theory gives rise to domain walls in the early Universe. While this is problematic if the walls are stable, it can be shown that a low energy ($T < 10^9\text{GeV}$) epoch dominated by transient domain walls can help to remove unwanted relics generic to supersymmetry. In eq.s (14) and (20) we obtain a requirement on the extent of secondary inflation caused by the domain walls, relating it to the energy scale $M_R$. We explored permissible values of $T_D$, (the temperature scale at which the domain walls finally disappear) for the walls to be unstable despite causing sufficient secondary inflation. The conditions appear as limits on differences of parameters in the supersymmetry breaking soft terms. The values in Table I and in eq. (30) can be used to constrain mechanisms of SUSY breaking and communication to the visible sector.

The sequence of cosmological events that take place in the two class of models is slightly different as indicated in table II. In both MSLRM and MSLR/P domain walls form when one of the two $\Omega$ fields spontaneously gets a vev at a scale $M_R \sim 10^6\text{GeV}$. The domains are distinguished by which of the two fields acquires a vev in them. The formation of domain walls results in their energy density dominating the energy density of the Universe, in turn causing a secondary inflation. This dilutes the gravitino and moduli fields. As the universe cools the triplet Higgs fields get a vev at a scale $M_{B-L} \sim 10^4\text{GeV}$. The parity breakdown
| Cosmology | Scale | Symmetry Group | MSLRP (GeV) | MSLRM (GeV) |
|-----------|-------|----------------|-------------|-------------|
| Ω or Ω_c get vev. | \( M_R \) | \( SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L} \) | 10^6 | 10^6 |
| Onset of wall dominated secondary inflation. | \( M_{B-L} \) | \( SU(3)_c \otimes SU(2)_L \otimes U(1)_R \otimes U(1)_{B-L} \) | 10^4 | 10^4 |
| Higgs triplet (\( \Delta' \)s) get vev | \( M_{B-L} \) | \( SU(3)_c \otimes SU(2)_L \otimes U(1)_Y \) (SUSY) | 10^3 | 10^3 |
| End of inflation and beginning of L-genesis | \( M_S \) | | | 10^3 |
| Wall disappearance temperature | \( T_D \) | | 10 − 10^3 | 10 − 10^2 |
| Secondary reheat temperature | \( T_R^* \) | | 10^3 − 10^4 | 10^3 |
| Electroweak breaking | \( M_W \) | \( SU(3)_c \otimes SU(2)_L \otimes U(1)_{Y} \) (non-SUSY) | 10^2 | 10^2 |
| Standard model | | \( SU(3)_c \otimes U(1)_{EW} \) | | |

TABLE II: Pattern of symmetry breaking and the slightly different sequence of associated cosmological events in the two classes of models

mechanism proposed in MSLRP, causes an asymmetry between \( L \) and \( R \) sectors at this stage. As such DW are de-stabilized and secondary inflation ends. The motion of the walls with a preferred direction of motion makes it possible for a mechanism for leptogenesis
proposed earlier [17] to be operative in MSLRP at this energy scale. In MSLRM however, parity is still unbroken. At a scale $M_s \sim 10^3$ GeV, SUSY breaking is mediated from the hidden sector to the visible sector in both the models. The soft terms which come into play at this stage break the parity in MSLRM explicitly. Thus, only at this stage do DW become unstable in MSLRM, thereby ending secondary inflation and commencing leptogenesis. The walls finally disappear in MSLRP at a temperature of $T_D \sim 10 - 10^3$ GeV unlike MSLRM where the same thing happens at a temperature range of $T_D \sim 10 - 10^2$ GeV. While the motion of the walls produces lepton asymmetry, the walls decay due to collisions dumping entropy into the medium and reheating the Universe. For MSLRP the reheate temperature from secondary inflation ($T_{R_s}$) can be estimated to range from $10^3 - 10^4$ GeV, whereas for MSLRM, it can be estimated to be $\lesssim 10^3$ GeV. Standard cosmology follows from this epoch onwards. The estimates of reheating are based on the energy stored in the walls and accord with the requirement that sphaleronic processes be effective in converting the lepton asymmetry generated by DW mechanism into baryon asymmetry. The values of reheating temperature also keep open the possibly of other TeV scale mechanisms for baryogenesis in these models.

In summary we find both MSLRM and MSLRP having salient cosmological features which can cure the problems of TeV scale supersymmetric theories. Both models bear further investigation due the constraints implied on the mechanism of supersymmetry breaking. The issue of obtaining non-thermal leptogenesis in this class of models is also an open question.

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APPENDIX A: F-FLATNESS AND D-FLATNESS CONDITIONS

The $F$-flatness conditions for the MSLRP are:

\[
F_{\Delta} = m_{\Delta} \Delta + a(\Delta \Omega - \frac{1}{2} \text{Tr} \Delta \Omega) = 0
\]

\[
F_{\Delta_c} = m_{\Delta} \Delta_c - a(\Delta_c \Omega_c - \frac{1}{2} \text{Tr} \Delta_c \Omega_c) = 0
\]

\[
F_{\Delta} = m_{\Delta} \Delta + i f L L^T \tau_2 + a(\Omega \Delta - \frac{1}{2} \text{Tr} \Omega \Delta) = 0
\]

\[
F_{\Delta_c} = m_{\Delta} \Delta_c + i f L_c L_c^T \tau_2 - a(\Omega_c \Delta_c - \frac{1}{2} \text{Tr} \Omega_c \Delta_c) = 0
\]

\[
F_{\Omega} = m_{\Omega} \Omega + a(\Delta \bar{\Delta} - \frac{1}{2} \text{Tr} \Delta \bar{\Delta}) = 0
\]

\[
F_{\Omega_c} = m_{\Omega} \Omega_c - a(\Delta_c \bar{\Delta}_c - \frac{1}{2} \text{Tr} \Delta_c \bar{\Delta}_c) = 0
\]

\[
F_{L} = 2 i f \tau_2 \Delta L = 0
\]

\[
F_{L_c} = 2 i f^* \tau_2 \Delta_c L_c = 0
\]  \quad (A1)

The $D$-flatness conditions are given by

\[
D_{Ri} = 2 \text{Tr} \Delta_i \tau_i \Delta_c + 2 \text{Tr} \bar{\Delta}_i \tau_i \bar{\Delta}_c + 2 \text{Tr} \bar{\Omega}_i \tau_i \Omega_c + L_c^i \tau_i L_c = 0
\]

\[
D_{Li} = 2 \text{Tr} \Delta^i \tau_i \Delta + 2 \text{Tr} \bar{\Delta}^i \tau_i \bar{\Delta} + 2 \text{Tr} \Omega^i \tau_i \Omega + L^i \tau_i L = 0
\]

\[
D_{B-L} = -L^i L + 2 \text{Tr} (\Delta^i \Delta - \bar{\Delta}^i \bar{\Delta}) + L_c^i L_c - 2 \text{Tr} (\Delta_c^i \Delta_c - \bar{\Delta}_c^i \bar{\Delta}_c) = 0
\]  \quad (A2)

[1] R. Kuchimanchi and R. N. Mohapatra, Phys. Rev. D 48, 4352 (1993).
[2] R. Kuchimanchi and R. N. Mohapatra, Phys. Rev. Lett. 75, 3989 (1995).
[3] M. Cvetič, Phys. Lett. B 164, 55 (1985).
[4] C. S. Aulakh, K. Benakli, and G. Senjanovic, Phys. Rev. Lett. 79, 2188 (1997), hep-ph/9703434.
[5] C. S. Aulakh, A. Melfo, A. Rasin, and G. Senjanovic, Phys. Rev. D 58, 115007 (1998), hep-ph/9712551.
[6] U. A. Yajnik and A. Sarkar, AIP Conf. Proc. 903, 685 (2007), hep-ph/0610161.
[7] T. W. B. Kibble, Phys. Rep. 67, 183 (1980).
[8] J. Ellis, J. E. Kim, and D. V. Nanopoulos, Phys. Lett. B 145, 181 (1984).
[9] D. H. Lyth and E. D. Stewart, Phys. Rev. D 53, 1784 (1996).
[10] T. Matsuda, Phys. Lett. B 486, 300 (2000).
[11] M. Kawasaki and F. Takahashi, Phys. Lett. B 618, 1 (2005), hep-ph/0410158.
[12] C. Grojean and G. Servant, Phys. Rev. D 75, 043507 (2007), hep-ph/0607107.
[13] W. Buchmuller, P. Di Bari, and M. Plumacher, Nucl. Phys. B 665, 445 (2003), hep-ph/0302092.
[14] W. Buchmuller, P. Di Bari, and M. Plumacher, New J. Phys. 6, 105 (2004), hep-ph/0406014.
[15] W. Fischler, G. F. Guidice, R. G. Leigh, and S. Paban, Phys. Lett. B 258, 45 (1991).
[16] N. Sahu and U. A. Yajnik, Phys. Rev. D 71, 023507 (2005).
[17] J. M. Cline, U. A. Yajnik, S. N. Nayak, and M. Rabikumar, Phys. Rev. D 66, 065001 (2002).
[18] L. Boubekeur (2002), hep-ph/0208003.
[19] T. Hambye, Nucl. Phys. B 633, 171 (2002).
[20] L. Boubekeur, T. Hambye, and G. Senjanovic, Phys. Rev. Lett. 93, 111601 (2004), hep-ph/0404038.
[21] E. J. Chun and S. Scopel, Phys. Lett. B 636, 278 (2006), hep-ph/0510170.
[22] S. Scopel, J. Phys. Conf. Ser. 39, 12 (2006).
[23] S. K. Majee, M. K. Parida, A. Raychaudhuri, and U. Sarkar, Phys. Rev. D 75, 075003 (2007), hep-ph/0701109.
[24] M. Joyce, T. Prokopec, and N. Turok, Phys. Rev. Lett. 75, 1695 (1995), hep-ph/9408339.
[25] M. Joyce, T. Prokopec, and N. Turok, Phys. Rev. D 53, 2958 (1996), hep-ph/9410282.
[26] J. M. Cline and K. Kainulainen, Phys. Rev. Lett. 85, 5519 (2000).
[27] J. M. Cline, M. Joyce, and K. Kainulainen, Phys. Lett. B 417, 79 (1998), 448, 321, (1999).
[28] D. Chang, R. N. Mohapatra, and M. K. Parida, Phys. Rev. D 30, 1052 (1984).
[29] D. Lindley, Mon. Not. Roy. Astron. Soc. 193, 593 (1980).
[30] C. L. Bennett and etal., Astrophys. J. Suppl. 148, 1 (2003), astro-ph/0302207.
[31] J. Preskill, S. P. Trivedi, F. Wilczek, and M. B. Wise, Nucl. Phys. B 363, 207 (1991).
[32] M. Dine and A. E. Nelson, Phys. Rev. D 48, 1277 (1993).