Testing Spontaneous Parity Violation at the LHC

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Abstract

We construct a supersymmetric $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model in which a discrete symmetry ($C$-parity) implements strict left-right symmetry in the scalar (Higgs) sector. Although two electroweak bidoublets are introduced to accommodate the observed fermion masses and mixings, a natural missing partner mechanism insures that a single pair of MSSM Higgs doublets survives below the left-right symmetry breaking scale. If this scale happens to lie in the TeV range, several new particles potentially much lighter than the $SU(2)_R$ charged gauge bosons $W^\pm_R$ will be accessible at the LHC.

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Left-right symmetric models, in which the observed parity violation arises due to spontaneous symmetry breaking, have been extensively studied in the past 1,2,3. Reference 3, in particular, considered supersymmetric (SUSY) models based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. It was shown in 3 how the MSSM $\mu$ problem can be neatly resolved and hybrid inflation realized. For the most part, the discussion in 3 assumed that the scalar (Higgs) sector of the theory did not respect left-right symmetry. A brief discussion near the end of the paper did, however, indicate that a left-right symmetric scalar sector may lead to several additional (non-MSSM) states in the low energy theory. Our main goal in this paper is to provide a realistic SUSY $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model in which strict left-right symmetry is enforced in all of the sectors by a discrete C-parity 4,5,6. (C-parity interchanges left and right and conjugates the representations. It also requires $g_L = g_R$, where $g_{L,R}$ denotes the $SU(2)_{L,R}$ gauge couplings respectively\(^1\).) We will assume that the scale of spontaneous parity violation ($M_R$) exceeds the SUSY breaking scale of around a TeV. In order to accommodate the observed fermion masses and mixings one needs, it turns out, two electroweak bidoublets. Fortunately, with manifest left-right symmetry in the scalar sector, a natural missing partner mechanism insures that with $M_R \gg M_W$, the low energy theory essentially coincides with the MSSM. However, if the scale of spontaneous parity violation $M_R$ happens to be in the multi-TeV range, the model predicts the existence of several new charged and neutral states which can be significantly lighter than, say $W_R$, and therefore accessible at the LHC. In the simplest models the spontaneous breaking of left-right symmetry can be achieved by employing $SU(2)_{L,R}$ doublet fields. Symmetry breaking with $SU(2)_{L,R}$ triplet fields is also possible and gives rise to additional states at low energy carrying two units of electric charge 7,8.

We extend, following 1, the Standard Model (SM) electroweak gauge group $SU(2)_L \times U(1)_Y$ with its chiral matter content to a left-right gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, where $B - L$ is baryon number minus lepton number and

$$Y = B - L + T_{3R},$$

(1)

where $T_{3R}$ is the third generator of $SU(2)_R$ in the Pauli basis. C-parity shows up as a $\mathbb{Z}_2$ automorphism of the gauge group interchanging $SU(2)_L$ with $SU(2)_R$ and conjugating

\(^{1}\) It is well known that C-parity (also known as D-parity) is contained in $SO(10)$ 4,5,6.
$U(1)_{B-L}$ (and also color). In such a model, matter turns up as the parity pairs $(3, 2, 1)^1_3 \oplus (\bar{3}, 1, 2)^{-1}_3$ and $(1, 2, 1)^{-1}_1 \oplus (1, 1, 2)^1_1$.

The electroweak symmetry breaking is implemented with a $C$-even bidoublet $(1, 2, 2)^0_0$. When it acquires a nonzero vacuum expectation value (VEV), it breaks both $SU(2)_L$ and $SU(2)_R$. However, this is not what we observe experimentally; we have not observed any $SU(2)_R$ gauge bosons so far. What this suggests is that $SU(2)_R \times U(1)_{B-L}$ – but not $SU(2)_L$ – is broken to $U(1)_Y$ at some energy scale higher than the electroweak scale – which might be the Grand Unified (GUT) scale or some other intermediate scale – and thereby breaking $C$-parity as well. If a Higgs field is responsible for this symmetry breaking, then it has to be a singlet under $SU(3)_C$ and $SU(2)_L$ but not under $SU(2)_R$.

The simplest choice for breaking $SU(2)_R \times U(1)_{B-L}$ down to $U(1)_Y$ would be an $SU(2)_R$ doublet $\Phi_R$. Later, we will also analyze $SU(2)_R$ triplets $\Delta_R$. Because of $C$-parity, we also need to include an $SU(2)_L$ doublet $\Phi_L$ as well. In a supersymmetric model, to cancel the gauge anomaly and to be able to construct simple superpotentials which can give rise to the Higgs mechanism, we also need to introduce the complex conjugates of these fields, $\Phi^c_L$ and $\Phi^c_R$ to the model. In addition, the conjugate fields are needed so that the $SU(2)_R \times U(1)_{B-L}$ breaking will not lead to some nonzero $D$-terms, thereby breaking SUSY at too high a scale. To summarize, it would seem that a minimal left-right model with $C$-parity requires at least three generations of matter, a Higgs bidoublet $H$ and possibly $\Phi_L$ and $\Phi_R$ (or left and right triplets) plus their $C$-parity conjugates.

Now from the fact that the $\Phi_R$’s and/or $\Delta_R$’s acquire nonzero VEVs but the $\Phi_L$’s and the $\Delta_L$’s do not, we see that the VEV of some superfields are not the same as their $C$-conjugates. This means that $C$-parity is spontaneously broken at the same scale as the $SU(2)_R$ breaking scale.

The minimal matter content consists of three generations of $L(1, 2, 1)^{-1}_1$, $L^c(1, 1, 2)^1_1$, $Q(3, 2, 1)^1_3$ and $Q^c(\bar{3}, 1, 2)^{-1}_3$. A Yukawa coupling like $LHL^c$ or $QHQ^c$ predicts that $Y^U = Y^D$ and $Y^E = Y^D$. This is a typical left-right prediction. However, we know that this is definitely not the case for all the three generations for the case of the quarks. One way around this relation is to double the number of bidoublets and make their electroweak VEVs point in different directions (i.e. the ratios $\langle H_{ui} \rangle / \langle H_{di} \rangle$ are different) and write down couplings to both of these bidoublets. We now need to explain why the VEVs point in different directions. Let us call the bidoublets $H_i, i = 1, 2$. The terms responsible for the
electroweak scale VEVs are the radiative corrections coming primarily from couplings to the third generation after SUSY is broken. To get MSSM at low energies, we only want a pair of $SU(2)_L$ doublets to survive at low energies. This can be achieved if say only the $H_u$ from $H_1$ and $H_d$ from $H_2$ remain light. This can be achieved via the “missing partner mechanism” \[9\]. Let us say we have the couplings $\Phi_L H_1 \Phi_R$ and $\Phi_L H_2 \Phi_R^c$. Then, after both $\Phi_R^c$ and $\Phi_R$ acquire $M$-scale VEVs, $\Phi_L$ pairs up with $H_d^1$ and $\Phi_R^c$ pairs up with $H_u^2$. The remaining light superfields are none other than $H_u^1$ and $H_d^2$. With such a mechanism, we can break the unwanted relation $Y^U = Y^D$ \[2\].

We note that even if we have the more general coupling
\[
W \supset c_{11} \Phi_L H_1 \Phi_R + c_{12} \Phi_L H_2 \Phi_R + c_{21} \Phi_L^c H_1 \Phi_R^c + c_{22} \Phi_L^c H_2 \Phi_R^c, \tag{2}
\]
with the coefficients free to vary, generically, unless $c_{11}c_{22} = c_{12}c_{21}$, we will still end up with the missing partner mechanism. The key here is to forbid the potentially huge $H_1 H_2$ coupling, which is taken care of automatically if we have $R$-symmetry. Without an $R$-symmetry, we need to come up with another mechanism to forbid such couplings together with nonrenormalizable couplings like $(\Phi_R^c \Phi_R)^n H_1 H_2 / \Lambda^{2n-1}$, where $\Lambda$ denotes some cutoff scale.

We may consider couplings like
\[
(L^c \Phi_R) (\Phi_R^c HL) / M^2, (L^c \Phi_R^c) (\Phi_R HL) / M^2, (Q^c \Phi_R) (\Phi_R^c HQ) / M^2, (Q^c \Phi_R^c) (\Phi_R HQ) / M^2 \tag{3}
\]
and their $C$-conjugates to break the left-right relation. These nonrenormalizable couplings can be gotten if we integrate over some massive fields with masses of order $M$, for instance. However, in a minimal model, these couplings will be suppressed by the Planck scale, which might lead to couplings which are too small.

Next, let us try to construct the Higgs sector step by step. Our first task is to come up with some mechanism to give the $\Phi_R$'s a nonzero VEV and one of the simplest ways of doing that is to introduce another $C$-even singlet superfield $S$ and write down the following superpotential \[11, 3, 10\]:
\[
W = \kappa S (\Phi_L^c \Phi_L + \Phi_R^c \Phi_R - M^2). \tag{4}
\]

Alternatively, with a single bidoublet one could introduce the superfields $(1, 3, 1)_0$ and $(1, 1, 3)_0$, and employ nonrenormalizable couplings involving the VEV $\langle(1, 1, 3)_0\rangle$ to break the relation $Y^U = Y^D$. \[2\]
The singlet $S$ plays an important role in hybrid inflation models \cite{11,3,10}. (For a recent analysis and additional references, see \cite{12}). Without any loss of generality, we may redefine the phase of $S$ and the $\Phi$’s so that both $\kappa$ and $M$ are real and positive\cite{3}.

We first note that in the absence of any gauge couplings, a generic superpotential of the form $Sf(\Phi)$ will consist mostly of moduli. The generic supersymmetric solution is $S = 0$ and $f = 0$. The latter condition singles out a complex submanifold of codimension 1 in $\Phi$ space. This moduli is not due to any symmetry, accidental or otherwise, in general but to holomorphy and the form of the superpotential. It may happen that in some cases, as in Eq. 4, that this result can also be reinterpreted in terms of an accidental $U(4)$ symmetry.

Since we only want SUSY to be broken at the TeV scale, to a good approximation, we can assume that the $D$ terms are zero at the scale of $M$. This forces $\phi_L$ and $\phi_R$, and $\phi^c_L$ and $\phi^c_R$ to acquire the same (complex conjugated) VEVs. We can always perform a gauge transformation so that these VEVs are identical and real. Since $SU(2)_L \times U(1)_Y$ is only broken at the electroweak scale, we need $\langle \phi_L \rangle = 0$.

Because of the moduli, the model as it currently stands contains some additional light fields. Because of this, we will not analyze it any further but will look into a more realistic model. As mentioned in the introduction, a realistic minimal model will likely contain two electroweak doublets with a missing partner mechanism. In addition, this pairing mechanism gets rid of precisely the light particles that \cite{3} had noticed. Because of this, we can get MSSM at low energies without the undesirable left-right mass relations.

Our model can be briefly summarized by the chiral superfield content of Table I and the

| superfield | representation | superfield | representation |
|------------|----------------|------------|----------------|
| $S$        | $(1, 1, 1)_0$  | $\Phi_R$   | $(1, 1, 2)_{-1}$ |
| $\Phi_L$   | $(1, 2, 1)_1$  | $\Phi^c_R$ | $(1, 1, 2)_1$   |
| $\Phi^c_L$ | $(1, 2, 1)_{-1}$ | $\Phi^c_R$ | $(1, 1, 2)_1$   |
| $H_1$      | $(1, 2, 2)_0$  |            |                |
| $H_2$      | $(1, 2, 2)_0$  |            |                |
| $Q_i$      | $(3, 2, 1)_{\frac{1}{4}}$ | $Q^c_i$   | $(\overline{3}, 1, 2)_{-\frac{1}{4}}$ |
| $L_i$      | $(1, 2, 1)_{-1}$ | $L^c_i$   | $(1, 1, 2)_1$   |

TABLE I: The chiral superfields in the minimal model.
TABLE II: The Higgs supermultiplet and gauge spectrum of the minimal model at the tree level in the limit where $M_3 \rightarrow 0$.

The superpotential

\[ W = \kappa S(\Phi_L^c \Phi_L + \Phi_R^c \Phi_R - M^2) + \alpha \Phi_L H_1 \Phi_R + \beta \Phi_L H_2 \Phi_R^c +
\]

\[ Y^U Q H_1 Q^c + Y^D Q H_2 Q^c + Y^E L H_2 L^c + Y^{Dirac} L H_1 L^c + \frac{1}{\Lambda} (\Phi_R \Phi_R L^c L^c + \Phi_L \Phi_L L L) \]

(5)

(This superpotential respects $Z_2$ matter parity.) Note the appearance of nonrenormalizable terms needed to provide Majorana masses to $\nu^c$ in $L^c$. The missing partner mechanism pairs up $\Phi_L$ with $H_{d1}$ and $\Phi_R^c$ with $H_{u2}$. The charged components of $\Phi_R$ and $\Phi_R^c$ are either eaten up or become sgoldstone partners of $W_R^\pm$, and some neutral components are eaten up or become sgoldstone partners of $Z_R$. The spectrum is summarized in Table II.

To get the seesaw mechanism to work, we need a $\Phi_R \Phi_R L^c L^c / \Lambda$ coupling. This can be obtained from a double seesaw mechanism if we introduce an $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ singlet neutrino superfield $N$ and the couplings $\Phi_R L^c N$ and $N^2$. It appears that any additional symmetry which commutes with all the other symmetries and allows all the couplings in Eq. 5 will also allow couplings like $S^2$ and $S^3$. However, such couplings will not change our analysis by much. Matter parity is still needed to forbid unwanted couplings.

Let us assume for the moment that our model has $U(1)_R$ symmetry and see where our analysis leads us. $S$ has an $R$-charge of 2. By $C$-parity, $\Phi_L$ and $\Phi_R$, $\Phi_L^c$ and $\Phi_R^c$, $Q$ and $Q^c$, and $L$ and $L^c$ will all have the same $R$-charge. $H_1$ and $H_2$ have to have the same $R$-charge, and $\Phi_L$ and $\Phi_L^c$ opposite $R$-charges. Putting all of this together, we conclude that
the $\mathcal{R}$-charges of the $\Phi$ fields have to be zero and the $\mathcal{R}$-charges of the bidoublets 2. The matter superfields have zero $\mathcal{R}$-charges. We can also immediately see that matter parity has to be a symmetry independent of $\mathcal{R}$-symmetry. This analysis nearly works except that the seesaw term would now have to break $\mathcal{R}$-symmetry. Because of this, we will not insist upon $\mathcal{R}$-symmetry.

On the other hand, if we really want to have $\mathcal{R}$-symmetry, we may consider using the mechanism proposed in Ref. [13]. Alternatively, instead of the seesaw mechanism, we may arrange following Ref. [14] to have a really tiny Dirac term. First, we must forbid the direct coupling $H_1 L^c L$ even while we require the $H_2 L^c L$ coupling, which certainly demands some explanation. Instead, we have a SUSY breaking sector involving a chiral superfield $X$ with an $R$-charge of 2 with an intermediate scale SUSY breaking $F_X \simeq M_{3/2}\Lambda$ and a Kähler term $X^\dagger H_1 L^c L/\Lambda$. This will induce a Dirac Yukawa coupling of order $M_{3/2}/\Lambda$, where $\Lambda$ is several orders of magnitude below the Planck scale.

We note that unlike other left-right models where $H_u$ and $H_d$ come from the same bidoublet, our model does not necessarily predict CP-violation because $H_u$ and $H_d$ come from different bidoublets and we can always perform a field redefinition to make both $\alpha$ and $\beta$ real and positive. Note however, that a tiny VEV for $H_{d1}$ and $H_{u2}$ will be induced. These VEVs are proportional to $M^{-2}$, which leads to negligible CP-violations provided that $M$ is large enough. The CP-violation coming from the matter Yukawa sector remains, as it must.

Next, we consider the case where $SU(2)_R \times U(1)_{B-L}$ is broken at some low energy scale. In Ref. [15], experimental bounds coming from CP-violation as measured by neutral kaon oscillations place a bound of about 2.5 TeV on the mass of the $W^\pm_R$ gauge boson. See also Ref. [16] which estimates a lower mass bound on $W^\pm_R$ of around 600 GeV and Ref. [17] which gives a lower bound of 1.6 TeV from the kaon mass splitting. From Table II with the dimensionless coefficients $\alpha, \beta$ of order unity or less, we predict the existence of several new particles which may be significantly lighter than $W^\pm_R$ and accessible at the LHC.

When the $SU(2)_R$ breaking scale is high, the VEVs of $H_{d1}$ and $H_{u2}$ go to zero and we can redefine the phases of $H_1$ and $H_2$ independently so that we do not have any contributions to CP-violation. However, as the breaking scale goes down, these VEVs now become larger and more significant and they contribute some amount to CP-violation from this sector.

We also note that the spontaneous breaking of $C$-parity gives rise to $\mathbb{Z}_2$ domain walls. If the symmetry breaking scale happens to lie below the reheating temperature, such domain
walls will appear after the universe has cooled down but they will not be inflated away. This is potentially problematic from a cosmological point of view.

There are at least two ways of giving $\nu^c$ (the neutral component of $L^c$) a mass when $M$ is of the order of a TeV or so; either through the double seesaw mechanism or via a $SU(2)_L$ triplet. The former case has already been discussed. We will now move on to the latter possibility, like that analyzed in Ref. [8]. With such a low scale mass for $\nu^c$, we need to suppress the Yukawa Dirac coupling by making it zero for instance, at the tree level. Radiative corrections will then generate a small value for the coupling resulting in the seesaw mechanism.

Let us now look at an alternative model with the $SU(2)_L$ triplets $\Delta_L(1,3,1)_2$, $\Delta^c_L(1,3,1)_{-2}$, $\Delta_R(1,1,3)_{-2}$ and $\Delta^c_R(1,1,3)_2$. We introduce the complex conjugates to cancel the gauge anomalies, among other things. We will still need to have Higgs bidoublets to give the matter Yukawa couplings at the renormalizable level. To break the left-right mass relation, we still need the missing partner mechanism which necessitates the presence of all the $\Phi$ fields\(^3\). Actually, since we are going to have so many additional particles above the $SU(2)_R$ breaking scale ($M$) anyway, in the case where it happens to be low, it might not matter so much if we have additional unpaired superfields at the supersymmetric level since their soft masses would be comparable to the masses that they would get from an $M$-scale pairing anyway. Of course, it is also possible to consider a model where $M$ happens to be large. But in that case, we need to find another model which reduces to MSSM. This will be the second $\Delta$ model that we will consider. In other words, this alternative model is really an extension of the model that we have been studying previously. Because of this, $SU(2)_R$ will be broken by both $\Phi_R$ as well as $\Delta_R$. The primary reason for introducing these triplets is to give $\nu^c$ a mass via the coupling $\Delta_R L^c L^c$.

Let us first consider the minimal (reduced) $\Delta$ model which gives rise to more low energy superfields than the MSSM. To break the up-down relation, we still need two Higgs bidoublets. This is more than what we have in MSSM, but all the additional Higgs fields can be made massive by the soft SUSY terms. If we assume that the SUSY breaking scale is smaller than the $SU(2)_R$ breaking scale, then at the LHC, we would expect to see a doubly charged Dirac fermion and two doubly charged scalars coming from $\Delta_R$ and $\Delta^c_R$; and a

\(^3\) In addition, without the $\Phi$’s, we will still be left with an unbroken $Z_{2R}$ gauge symmetry.
doubly charged Dirac fermion and two doubly charged scalars coming from $\Delta_L$ and $\Delta^c_L$, a charged Dirac fermion and two charged scalars coming from $\Delta_L$ and $\Delta^c_L$, a neutral Dirac and four neutral scalars coming from $\Delta_L$ and $\Delta^c_L$, and two charged Dirac Higgsinos and two neutral Dirac Higgsinos and three charged Higgs and seven neutral Higgs coming from the two bidoublets.

Let us now turn to the second $\Delta$ model which reduces to MSSM if the left-right symmetry breaking scale $M \gg \text{TeV}$. While we will be primarily interested in the case where $M$ is low, as this may give rise to new physics at the LHC, this model will still be acceptable if $M$ is high. A renormalizable superpotential of the form

$$ W \supset \kappa S(\Phi_L^c \Phi_L + \Phi_R^c \Phi_R + \rho(\Delta_L^c \Delta_L + \Delta_R^c \Delta_R) - M^2) + M'(\Delta_L^c \Delta_L + \Delta_R^c \Delta_R) + \alpha(\Delta_R \Phi_R^c \Phi_R + \Delta_L \Phi_L^c \Phi_L) + \beta(\Delta_R \Phi_R^c \Phi_R + \Delta_L \Phi_L^c \Phi_L) $$

(6)

will give rise to nonzero VEVs for both the $\Phi$’s and the $\Delta$’s. The equation $F_{\Delta_R} = F_{\Delta_R^c} = 0$ causes $\langle \Delta \rangle$ to be proportional to $\langle \Phi \rangle^2$. $F_S = 0$ causes the VEVs to be nonzero. Without any loss of generality, we may assume that there is no $\Phi_L^c \Phi_L + \Phi_R^c \Phi_R$ term because any such term can always be reabsorbed into a field redefinition of $S$ by some shift.

To pair up the $\Delta_L$’s, we can introduce $K_L(3,2)_1$, $K_R(2,3)_1$, $K_L^c(3,2)_1$ and $K_R^c(2,3)_1$. Introduce the renormalizable couplings $\Delta_L K_L^c \Phi_R$, $\Delta_R K_R \Phi_L$, $\Delta_L^c K_L \Phi_R^c$ and $\Delta_R^c K_R^c \Phi_R$ and


doubly charged Dirac fermion and two doubly charged scalars coming from $\Delta_L$ and $\Delta^c_L$, a charged Dirac fermion and two charged scalars coming from $\Delta_L$ and $\Delta^c_L$, a neutral Dirac and four neutral scalars coming from $\Delta_L$ and $\Delta^c_L$, and two charged Dirac Higgsinos and two neutral Dirac Higgsinos and three charged Higgs and seven neutral Higgs coming from the two bidoublets.

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$$ W \supset \kappa S(\Phi_L^c \Phi_L + \Phi_R^c \Phi_R + \rho(\Delta_L^c \Delta_L + \Delta_R^c \Delta_R) - M^2) + M'(\Delta_L^c \Delta_L + \Delta_R^c \Delta_R) + \alpha(\Delta_R \Phi_R^c \Phi_R + \Delta_L \Phi_L^c \Phi_L) + \beta(\Delta_R \Phi_R^c \Phi_R + \Delta_L \Phi_L^c \Phi_L) $$

(6)

will give rise to nonzero VEVs for both the $\Phi$’s and the $\Delta$’s. The equation $F_{\Delta_R} = F_{\Delta_R^c} = 0$ causes $\langle \Delta \rangle$ to be proportional to $\langle \Phi \rangle^2$. $F_S = 0$ causes the VEVs to be nonzero. Without any loss of generality, we may assume that there is no $\Phi_L^c \Phi_L + \Phi_R^c \Phi_R$ term because any such term can always be reabsorbed into a field redefinition of $S$ by some shift.

To pair up the $\Delta_L$’s, we can introduce $K_L(3,2)_1$, $K_R(2,3)_1$, $K_L^c(3,2)_1$ and $K_R^c(2,3)_1$. Introduce the renormalizable couplings $\Delta_L K_L^c \Phi_R$, $\Delta_R K_R \Phi_L$, $\Delta_L^c K_L \Phi_R^c$ and $\Delta_R^c K_R^c \Phi_R$ and

\[ \text{TABLE III: The chiral superfield content of the reduced $\Delta$ model.} \]
also the couplings \( K_L K_L \Delta_R^c, K_L^c K_R^c \Delta_R \) and their \( \mathcal{C} \)-conjugates. The \( \Delta_L \)'s pair up and the \( 3_0 \)'s of the \( K \)'s also pair up and all the masses are of the \( \Lambda_R \) scale. A coupling like \( K_L K_L^c \) would be disastrous because the \( \Delta_L \)'s will only get seesaw contributions to their masses. The full chiral superfield content of this model is given in Table IV.

We also need to give masses to the doubly-charged component of \( \Delta_R \). As already mentioned, a direct \( \Delta_R^c \Delta_R \) will not do. So, let us introduce an additional \( A_L(3, 1)_0 \) and \( A_R(1, 3)_0 \) and the couplings \( \Delta_L \Delta^c_L A_L \) and \( \Delta_R \Delta^c_R A_R \). The last term in addition to the mass term \( A_R^2 \) (and its \( \mathcal{C} \)-conjugate) will induce a nonzero VEV \( \sim M \) for \( A_R \) since \( \Delta_R \) and \( \Delta_R^c \) already have nonzero VEVs. With so many additional light particles at the TeV scale, we will have plenty of new physics to work with at the LHC.
In summary, the superpotential will contain the following renormalizable terms:

\[ W \supset S, S (\Phi^c_L \Phi_L + \Phi^c_R \Phi_R), (\Delta^c_L \Delta_L + \Delta^c_R \Delta_R), \]

\[ (\Delta^c_L \Phi_L \Phi^c_L + \Delta^c_R \Phi^c_R \Phi^c_R), (\Delta^c_L \Phi_L \Phi^c_L + \Delta^c_R \Phi_R \Phi^c_R), (A_L \Phi^c_L \Phi_L + A_R \Phi^c_R \Phi_R), A^2_L, A^2_R, \]

\[ H_1 \Phi_L \Phi_R, H_2 \Phi^c_L \Phi^c_R, \]

\[ K_L \Delta_L \Phi_R + K_R \Delta_R \Phi_L, K^c_L \Delta^c_L \Phi_R + K^c_R \Delta^c_R \Phi_L, K^c_L K_L, K^c_R K_R. \]  \hspace{1cm} (7)

To sum up, we predict several \( SU(2)_L \) Higgs triplets, \( SU(2)_L \) Higgs doublets, and \( SU(2)_L \) singlets at low energies. The \( \Delta_L \)'s and both \( K_L \) and \( K_R \) will contribute doubly charged particles. In particular, we will have six doubly charged scalars and three doubly charged Dirac particles. These particles all have masses comparable to or smaller than the mass of \( W^\pm_R \).

I. CONCLUSION

In this paper we have insisted that the scalar (Higgs) sector of models based on symmetry groups such as \( SU(2)_L \times SU(2)_R \times U(1)_{B-L} \) should respect \( C \)-parity. We have considered a variety of models to show how the MSSM can be recovered at energies below the left-right \( (C \)-parity) symmetry breaking scale. If the latter happens to lie in the TeV range a plethora of new particles should be accessible at the LHC.

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