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Spontaneous parity violation and minimal Higgs models

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Abstract. In this paper we present a model for the spontaneous breaking of parity with two Higgs doublets and two neutral Higgs singlets which are even and odd under $\mathcal{D}$-parity. The condition $\upsilon_{\mathrm{R}} \gg \upsilon_{\mathrm{L}}$ can be satisfied without introducing bidoublets, and it is induced by the breaking of $\mathcal{D}$-parity through the vacuum expectation value of the odd Higgs singlet. Examples of left–right symmetric and mirror fermions models in grand unified theories are presented.

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1 Introduction

Left–right symmetric models with spontaneous parity breaking offer a natural explanation for the parity asymmetry observed in nature. The gauge group $SU(2)_{\mathrm{L}} \otimes SU(2)_{\mathrm{R}} \otimes U(1)_{B-L}$ fixes the interactions and generalizes the standard electroweak theory. However the fundamental fermionic representation and the Higgs sector are not completely determined. For the fermions one can have two possibilities: new right-handed doublets, as in the earlier models [1, 2], or new mirror fermions [3–6]. The Higgs sector has more possibilities and introduces more unknown parameters in the model. It is highly desirable to have the minimum number of unknown parameters in order to compare models and experimental data. A recent work in this direction was done by Brahmachari, Ma and Sarkar [7]. With two Higgs doublets, χ_{R} and χ_{L}, and a dimension five operator they have proposed a left–right model including fermion masses. Another mirror model with two Higgs doublets and two Higgs singlets was developed in [4–6]. In both cases the condition

$$\upsilon_{\mathrm{R}} \gg \upsilon_{\mathrm{L}} \tag{1}$$

must be satisfied. The present experimental bound is $\upsilon_{\mathrm{R}} > 30\upsilon_{\mathrm{L}}$ [4–6]. Later on, Siringo [8] revived an earlier remark by Senjanovic and Mohapatra [9] that the above condition cannot be satisfied in models with only two Higgs doublets. These remarks seem to leave open the possibility that only scalar bidoublets could break parity in a consistent way. This possibility exists but has the unpleasant feature of a large number of Higgs fields and undetermined parameters.

However, there is other elegant way [10, 11] to produce the condition (1) by introducing one singlet Higgs which is odd under $\mathcal{D}$-parity. The difference with the $\mathcal{P}$-parity breaking is that in $\mathcal{D}$-parity the vacuum expectation value (VEV) of a parity odd field can be spontaneously broken without breaking the left–right symmetry. In consequence, the gauge coupling constants of $SU(2)_{\mathrm{R}}$ and $SU(2)_{\mathrm{L}}$ also can be different and left-handed and right-handed scalars can have different masses and VEVs.

In the present paper, we extend the previous analysis to include also a singlet Higgs field which is even under $\mathcal{D}$-parity and mixes with the odd field. We show that this mixing term also contributes to the hierarchy relation (1). We include in our study two examples of grand unified theories where Higgs singlets transforming under $\mathcal{D}$-parity are assigned to the L–R symmetric model and to a mirror fermion model.

2 The scalar potential and the breaking of the L–R symmetry

There are two forms of breaking parity spontaneously: the first is to identify the discrete symmetry Z_2 that interchanges the groups $SU(2)_{\mathrm{L}}$ and $SU(2)_{\mathrm{R}}$ of the Lorentz group $O(3,1)$ as the parity operator $\mathcal{P}$ that transforms the Higgs bosons $\chi_{\mathrm{L}} \overset{\mathcal{P}}{\longleftrightarrow} \chi_{\mathrm{R}}$ and also $W_{\mathrm{L}} \overset{\mathcal{P}}{\longleftrightarrow} W_{\mathrm{R}}$. So, when $SU(2)_{\mathrm{R}}$ is broken in the symmetric L–R model, parity $\mathcal{P}$ is also broken. The second possibility of spontaneously breaking the parity symmetry is through the VEV of an odd scalar field which preserves L–R symmetry. This type of parity is called $\mathcal{D}$-parity; it is a generator of larger groups that contain the product $SU(2)_{\mathrm{L}} \otimes SU(2)_{\mathrm{R}}$ as a subgroup. This second possibility is very interesting because it allows for $\langle\chi_{\mathrm{L}}\rangle \ll \langle\chi_{\mathrm{R}}\rangle$ with different coupling constants for

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$SU(2)_{\rm L}$ and $SU(2)_{\rm R}$ and different masses for the Higgs fields.

Our model for the scalar potential includes two doublets and two singlets Higgs fields. These singlets and doublets transform under $\mathcal{D}$-parity as $S_M \overset{\mathcal{D}}{\longleftrightarrow} S_M$, $S_D \overset{\mathcal{D}}{\longleftrightarrow} -S_D$ and $\chi_{\rm L} \overset{\mathcal{D}}{\longleftrightarrow} \chi_{\rm R}$, if in the model there are no $\mathcal{CP}$ violating terms or no complex Yukawa couplings. We propose the following invariant potential under $G_{3221} = SU(3)_C \otimes SU(2)_{\rm L} \otimes SU(2)_{\rm R} \otimes U(1)_{B-L}$ for the Higgs fields:

$$\begin{aligned} V(\chi_{\rm L},\chi_{\rm R},S_D,S_M) = {} & \mu^2\left(\chi_{\rm L}^\dagger\chi_{\rm L}+\chi_{\rm R}^\dagger\chi_{\rm R}\right) \\ & -\lambda_\chi\left(\chi_{\rm L}^\dagger\chi_{\rm L}+\chi_{\rm R}^\dagger\chi_{\rm R}\right)^2 \\ & -m_D^2S_D^2-\eta_DS_D^3-\lambda_DS_D^4 \\ & -m_M^2S_M^2-\eta_MS_M^3-\lambda_MS_M^4 \\ & +M_DS_D\left(\chi_{\rm R}^\dagger\chi_{\rm R}-\chi_{\rm L}^\dagger\chi_{\rm L}\right) \\ & +M_MS_M\left(\chi_{\rm L}^\dagger\chi_{\rm L}+\chi_{\rm R}^\dagger\chi_{\rm R}\right) \\ & +\lambda S_DS_M\left(\chi_{\rm R}^\dagger\chi_{\rm R}-\chi_{\rm L}^\dagger\chi_{\rm L}\right) \\ & +\left(\varepsilon_DS_D^2+\varepsilon_MS_M^2\right)\left(\chi_{\rm L}^\dagger\chi_{\rm L}+\chi_{\rm R}^\dagger\chi_{\rm R}\right) \\ & -\kappa\left(\left(\chi_{\rm L}^4\right)^\dagger+\chi_{\rm L}^4+\left(\chi_{\rm R}^4\right)^\dagger+\chi_{\rm R}^4\right). \end{aligned} \tag{2}$$

Our motivation for taking this potential is the fact that S_M and S_D do not necessarily belong to the same irreducible multiplet of Higgs fields. In consequence, it is also possible that these fields are mixed. If this is the case, when $\langle S_D\rangle = s_D$ and $\langle S_M\rangle = s_M$ the potential terms that contribute to the masses of Higgs fields $\chi_{\rm L}$ and $\chi_{\rm R}$ are

$$\begin{aligned} V_{\rm mass}(\chi_{\rm L},\chi_{\rm R}) = {} & \left(\mu^2+\varepsilon_Ds_D^2+\varepsilon_Ms_M^2+M_Ms_M\right) \\ & \times\left(|\chi_{\rm L}|^2+|\chi_{\rm R}|^2\right) \\ & +\left(M_Ds_D+\lambda s_Ds_M\right)\left(|\chi_{\rm R}|^2-|\chi_{\rm L}|^2\right), \end{aligned} \tag{3}$$

from which we obtain the masses

$$m_{\rm R}^2 = \mu^2+\varepsilon_Ds_D^2+\varepsilon_Ms_M^2+M_Ms_M+M_Ds_D+\lambda s_Ds_M\,, \tag{4}$$

$$m_{\rm L}^2 = \mu^2+\varepsilon_Ds_D^2+\varepsilon_Ms_M^2+M_Ms_M-M_Ds_D-\lambda s_Ds_M\,. \tag{5}$$

Now we impose the hierarchy condition in the previous equations such that $m_{\rm R}^2 \ll s_D^2 \ll s_M^2$. It is necessary to indicate that $\upsilon_{\rm L}$ breaks the electroweak symmetry and $\upsilon_{\rm R}$ breaks the L–R symmetry close to the TeV scale. So we can have, for example, $\langle\chi_{\rm L}\rangle = \upsilon_{\rm L} \sim m_{\rm L} \sim 100\,\text{GeV}$ and $\langle\chi_{\rm R}\rangle = \upsilon_{\rm R} \sim m_{\rm R} \sim 10\,\text{TeV} \gg \upsilon_{\rm L}$. It also must be noted that if S_D and S_M are in the same multiplet, then several possible mixing terms in the potential possibility are absent.

Let us now suppose that there are no $\mathcal{CP}$ violating terms and that all VEVs are considered to be real. With $\langle\chi_{\rm L}\rangle = \begin{pmatrix}0\\ \upsilon_{\rm L}\end{pmatrix}$, $\langle\chi_{\rm R}\rangle = \begin{pmatrix}0\\ \upsilon_{\rm R}\end{pmatrix}$, it is possible to show that the minimum conditions for the potential are given by

$$\begin{aligned} \frac{\partial V}{\partial \upsilon_{\rm L}} &= 2\upsilon_{\rm L}\left[\mu^2-2\lambda_\chi\left(\upsilon_{\rm L}^2+\upsilon_{\rm R}^2\right)-M_Ds_D+M_Ms_M\right. \\ &\qquad \left.-\lambda s_Ds_M+\varepsilon_Ds_D^2+\varepsilon_Ms_M^2-4\kappa\upsilon_{\rm L}^2\right] \\ &= 0\,, \end{aligned} \tag{6}$$

$$\begin{aligned} \frac{\partial V}{\partial \upsilon_{\rm R}} &= 2\upsilon_{\rm R}\left[\mu^2-2\lambda_\chi\left(\upsilon_{\rm L}^2+\upsilon_{\rm R}^2\right)+M_Ds_D+M_Ms_M\right. \\ &\qquad \left.+\lambda s_Ds_M+\varepsilon_Ds_D^2+\varepsilon_Ms_M^2-4\kappa\upsilon_{\rm R}^2\right] \\ &= 0\,. \end{aligned} \tag{7}$$

From these equations we have

$$\begin{aligned} \upsilon_{\rm L}\frac{\partial V}{\partial \upsilon_{\rm R}}-\upsilon_{\rm R}\frac{\partial V}{\partial \upsilon_{\rm L}} &= 4\upsilon_{\rm L}\upsilon_{\rm R}\left[M_Ds_D+\lambda s_Ds_M\right. \\ &\qquad \left.-2\kappa\left(\upsilon_{\rm R}^2-\upsilon_{\rm L}^2\right)\right] \\ &= 0\,. \end{aligned} \tag{8}$$

Now we require non-trivial solutions such that $\upsilon_{\rm L} \neq \upsilon_{\rm R} \neq 0$. Thus we obtain the desired hierarchy:

$$\upsilon_{\rm R}^2-\upsilon_{\rm L}^2 = \frac{s_D(M_D+\lambda s_M)}{2\kappa}\,. \tag{9}$$

An important point to be noted in the previous equation is that the breaking effect due to the singlet S_M is sub-dominant with respect to S_D that breaks $\mathcal{D}$-parity when $\langle S_D\rangle = s_D$. Additionally, if $s_D = 0$ then $\mathcal{D}$-parity is conserved and the L–R symmetry condition is recovered, $\upsilon_{\rm R} = \upsilon_{\rm L}$. We have shown that in our potential it is possible to construct models with L–R symmetry and to produce an hierarchy between the breaking scale of $SU(2)_{\rm R}$ and the electroweak scale. The main ingredient is the presence of two Higgs singlets to generate the minimum of the potential. The crucial point in this sense is the inclusion of the mixing term $\lambda S_DS_M(\chi_{\rm R}^\dagger\chi_{\rm R}-\chi_{\rm L}^\dagger\chi_{\rm L})$, which is possible if S_M and S_D belong to different irreducible representations. As in the previous term, also the term $M_DS_D(\chi_{\rm R}^\dagger\chi_{\rm R}-\chi_{\rm L}^\dagger\chi_{\rm L})$ breaks the L–R symmetry. It is also fundamental to fine-tune the parameters of the model at the tree level in order to assure that $\upsilon_{\rm R}$ does not destabilize the $\upsilon_{\rm L}$ value. Thus, from (5)–(7) we have

$$m_{\rm L}^2-2(\lambda_\chi+2\kappa)\upsilon_{\rm L}^2 = 2\lambda_\chi\upsilon_{\rm R}^2\,. \tag{10}$$

3 An $SO(10)$ L–R symmetric model

There is known a GUT context in which one can embed the L–R symmetric model; this is based on $SO(10)$ through its maximal subgroup as in the Pati–Salam [1, 2] approach: $G_{\rm PS} = SU(4)_C \otimes SU(2)_{\rm L} \otimes SU(2)_{\rm R}$. The idea consists of breaking $\mathcal{D}$-parity below the breaking of $SO(10)$ as shown in the following breaking chain:

$$\begin{aligned} SO(10) &\xrightarrow{S_M} G_{\rm PS}\otimes\mathcal{D} \\ &\xrightarrow{S_D} SU(3)_C\otimes SU(2)_{\rm L}\otimes SU(2)_{\rm R}\otimes U(1)_{B-L} \\ &\xrightarrow{\chi_{\rm R}} SU(3)_C\otimes SU(2)_{\rm L}\otimes U(1)_Y \\ &\xrightarrow{\chi_{\rm L}} SU(3)_C\otimes U(1)_{\rm em}\,. \end{aligned} \tag{11}$$

Table 1. Higgs representations for the breaking chain (11)

$S_M \sim$	$\{\mathbf{54}\}$	$\supset [\mathbf{1},\mathbf{1},\mathbf{1}]$	$\sim (\mathbf{1},\mathbf{1},\mathbf{1},0)$
$S_D \sim$	$\{\mathbf{45}\}$	$\supset [\mathbf{15},\mathbf{1},\mathbf{1}]$	$\supset (\mathbf{1},\mathbf{1},\mathbf{1},0)$
$\chi_R \sim$	$\{\mathbf{144}^*\}$	$\supset [\mathbf{4},\mathbf{1},\mathbf{2}]$	$\supset (\mathbf{1},\mathbf{1},\mathbf{2},-1)$
$\chi_L \sim$	$\{\mathbf{144}\}$	$\supset [\mathbf{4},\mathbf{2},\mathbf{1}]$	$\supset (\mathbf{1},\mathbf{2},\mathbf{1},-1)$

The quantum numbers for the Higgs representations that produce the pattern (11) are given in Table 1.

Our notation is as follows: the representations between { } correspond to $SO(10)$, those with [ ] correspond to $G_{\rm PS}$, and those with ( ) correspond to $G_{3221} = SU(3)_C \otimes SU(2)_{\rm L} \otimes SU(2)_{\rm R} \otimes U(1)_{B-L}$. In our model a point different from the approach of [12] is that we are using the singlet component $[\mathbf{1},\mathbf{1},\mathbf{1}]$ of $\{\mathbf{54}\}$ in order to break $SO(10)$ down to $G_{\rm PS} \otimes \mathcal{D}$; because this is a symmetric representation, it is $\mathcal{D}$-even [10,11,13], which is different from the $[\mathbf{1},\mathbf{1},\mathbf{1}]$ component of $\{\mathbf{210}\}$, which is $\mathcal{D}$-odd, as required in our analysis in the previous section. Note also that the neutral component $(\mathbf{1},\mathbf{1},\mathbf{1},0) \subset [\mathbf{15},\mathbf{1},\mathbf{1}] \subset \{\mathbf{45}\}$ of $SO(10)$ is $\mathcal{D}$-odd under G_{3221}. Thus, we expect that the VEV of $S_D \sim \{\mathbf{45}\}$, which will induce the breaking of the L–R symmetry, will depend of the VEV of S_M. This choice could allow the breaking of the L–R symmetry to occur close to the electroweak scale; let us say on the scale of a few TeV.

The G_{3221} invariant Higgs potential of (2) could come from the following $SO(10)$ potential:

$$\begin{aligned}\mathcal{L} = {} & \mu^2(\mathbf{144}^* \times \mathbf{144}) + \lambda_\chi(\mathbf{144}^* \times \mathbf{144})^2 \\ & + m_D^2(\mathbf{45})^2 + \eta_D(\mathbf{45})^3 + \lambda_D(\mathbf{45})^4 \\ & + m_M^2(\mathbf{54})^2 + \eta_M(\mathbf{54})^3 + \lambda_M(\mathbf{54})^4 \\ & + M_D(\mathbf{45})(\mathbf{144}^* \times \mathbf{144}) + M_M(\mathbf{54})(\mathbf{144}^* \times \mathbf{144}) \\ & + \lambda(\mathbf{54} \times \mathbf{45})(\mathbf{144}^* \times \mathbf{144}) \\ & + \left(\varepsilon_D(\mathbf{45})^2 + \varepsilon_M(\mathbf{54})^2\right)(\mathbf{144}^* \times \mathbf{144}) \\ & + \kappa[(\mathbf{144}^*)^4 + (\mathbf{144})^4]\,. \end{aligned} \tag{12}$$

Let us notice that the term $(\mathbf{54} \times \mathbf{45})(\mathbf{144}^* \times \mathbf{144})$ is possible if the interactions between $(\mathbf{144}^* \times \mathbf{144})$ and $(\mathbf{54} \times \mathbf{45})$ are mediated by the gauge boson in the $\{\mathbf{45}\}$ or $\{\mathbf{54}\}$ representations.

The corresponding hypercharges are given by

$$\frac{Y}{2} = T_{3\rm R} + \frac{B-L}{2}\,. \tag{13}$$

The left-handed ordinary fermions are contained in

$$\{\mathbf{16}\}_{i\rm L} = \overbrace{q_{\rm L} = \begin{pmatrix} u \\ d \end{pmatrix}_{\rm L} (\mathbf{3},\mathbf{2},\mathbf{1},1/3) \oplus l_{\rm L} = \begin{pmatrix} \nu \\ e \end{pmatrix}_{\rm L} (\mathbf{1},\mathbf{2},\mathbf{1},-1)}^{[\mathbf{4},\mathbf{2},\mathbf{1}]}$$
$$\oplus \overbrace{q_{\rm L}^C = \begin{pmatrix} d^C \\ u^C \end{pmatrix}_{\rm L} (\mathbf{3}^*,\mathbf{1},\mathbf{2},-1/3) \oplus l_{\rm L}^C = \begin{pmatrix} e^C \\ \nu^C \end{pmatrix}_{\rm L} (\mathbf{1},\mathbf{1},\mathbf{2},1)}^{[\mathbf{4}^*,\mathbf{1},\mathbf{2}]}\,. \tag{14}$$

The Majorana and Dirac masses in our model can arise from dimension five effective operators such as

$$\begin{aligned} \mathcal{O}_1 &= \frac{1}{\Lambda_Q}(\overline{q_{\rm L}}\chi_{\rm L})(\overline{q_{\rm R}}\chi_{\rm R}^*)\,, & \mathcal{O}_2 &= \frac{1}{\Lambda_Q}(\overline{q_{\rm L}}\chi_{\rm L}^*)(q_{\rm R}\chi_{\rm R})\,, \\ \mathcal{O}_3 &= \frac{1}{\Lambda_D}(\overline{l_{\rm L}}\chi_{\rm L})(l_{\rm R}\chi_{\rm R}^*)\,, & \mathcal{O}_4 &= \frac{1}{\Lambda_D}(\overline{l_{\rm L}}\chi_{\rm L}^*)(l_{\rm R}\chi_{\rm R})\,, \\ \mathcal{O}_5 &= \frac{1}{\Lambda_M}(l_{\rm L}\chi_{\rm L}^*)(l_{\rm L}\chi_{\rm L}^*)\,, & \mathcal{O}_6 &= \frac{1}{\Lambda_M}(l_{\rm R}\chi_{\rm R}^*)(l_{\rm R}\chi_{\rm R}^*)\,, \end{aligned}$$

where $\Lambda_{Q,D,M}$ are the masses in the GUT scale. The first two operators $\mathcal{O}_1$ and $\mathcal{O}_2$ give the quark masses from which the existence is assumed of the matter fields $\mathcal{Q}_{\rm L,R} \sim (\mathbf{3},\mathbf{1},\mathbf{1},4/3) \subset (\mathbf{15},\mathbf{1},\mathbf{1}) \subset \{\mathbf{45}\}$ or $\{\mathbf{120}\}$ in order to generate the operator $\mathcal{O}_1$ and $\mathcal{M}_{\rm L,R} \sim (\mathbf{3},\mathbf{1},\mathbf{1},-2/3) \subset (\mathbf{6},\mathbf{1},\mathbf{1}) \subset \{\mathbf{10}\}$, $\{\mathbf{126}\}$ or $\mathcal{M}_{\rm L,R} \sim (\mathbf{3},\mathbf{1},\mathbf{1},-2/3) \subset (\mathbf{10},\mathbf{1},\mathbf{1}) \subset \{\mathbf{120}\}$ in order to generate the operator $\mathcal{O}_2$. On the other side, the operators $\mathcal{O}_3$ and $\mathcal{O}_4$ can be generated by the fermionic matter fields $\mathcal{P}_{\rm L,R} \sim (\mathbf{1},\mathbf{1},\mathbf{1},0) \subset (\mathbf{1},\mathbf{1},\mathbf{1}) \subset \{\mathbf{54}\}$ and $\mathcal{S}_{\rm L,R} \sim (\mathbf{1},\mathbf{1},\mathbf{1},-2) \subset (\overline{\mathbf{10}},\mathbf{1},\mathbf{1}) \subset \{\mathbf{120}\}$. This is shown in Fig. 1.

To obtain the operators $\mathcal{O}_5$ and $\mathcal{O}_6$ let us observe that the operator $(\{\mathbf{16}\}_{\rm L}\{\mathbf{16}\}_{\rm L})(\{\mathbf{144}^*\}\{\mathbf{144}^*\})$ can be obtained through the mediation of the fermions in the $\{\mathbf{45}\}$ and $\{\mathbf{210}\}$ representations as it is shown in Fig. 2. Operators of dimension five of the Majorana type $\mathcal{O}_5$ can be obtained from Fig. 2a, where the inclusion is necessary of the fermionic matter term $\{\mathbf{45}\}$ in the $(\mathbf{15},\mathbf{1},\mathbf{1})$ component. This contribution to the neutrino mass is large, because this fermion term corresponds to the GUT scale. To implement the see-saw mechanism we need to consider a term as shown in Fig. 2b which includes the $\mathcal{D}$-parity effect through the Higgs singlet $\{\mathbf{45}\}$ and also a term that preserves it through the Higgs singlet $\{\mathbf{54}\}$. Details of the calculation and their relevance for generating magnetic moments for the neutrinos and charged leptons will be presented elsewhere [14].

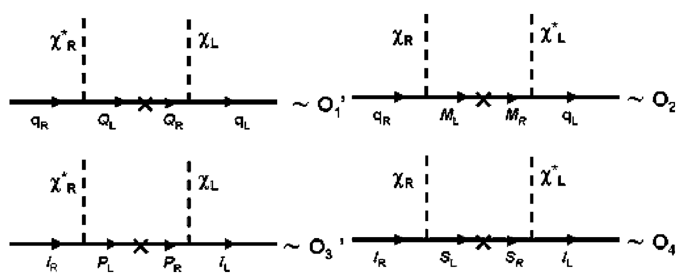

Fig. 1. Diagrams producing the $\mathcal{O}_{1-4}$ operators of dimension five

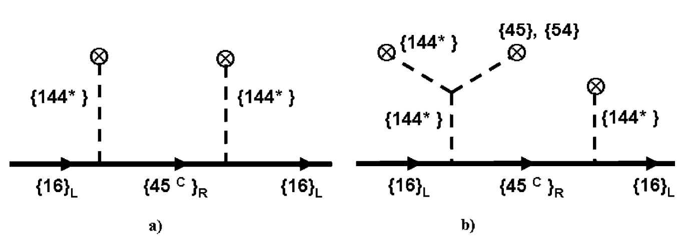

Fig. 2. Diagrams for generating neutrino masses

Table 2. Higgs representations for the breaking chain (15)

$S_M \sim$	$\{\mathbf{210}\}$	$\supset [\mathbf{1},\mathbf{1},\mathbf{1}]$	$\sim (\mathbf{1},\mathbf{1},\mathbf{1},0)$
$S_D \sim$	$\{\mathbf{210}\}$	$\supset [\mathbf{15},\mathbf{1},\mathbf{1}]$	$\supset (\mathbf{1},\mathbf{1},\mathbf{1},0)$
$\chi_R \sim$	$\{\mathbf{16}^*\}$	$\supset [\mathbf{4},\mathbf{1},\mathbf{2}]$	$\supset (\mathbf{1},\mathbf{1},\mathbf{2},-1)$
$\chi_L \sim$	$\{\mathbf{16}\}$	$\supset [\mathbf{4},\mathbf{2},\mathbf{1}]$	$\supset (\mathbf{1},\mathbf{2},\mathbf{1},-1)$

We have another possibility to embed $SU(2)_{\mathrm{L}} \otimes SU(2)_{\mathrm{R}} \otimes U(1)_{B-L}$ based in the breaking chain [15]

$$\begin{aligned} SO(10) &\xrightarrow{S_M} G_{\mathrm{PS}} \\ &\xrightarrow{S_D} SU(3)_C \otimes SU(2)_{\mathrm{L}} \otimes SU(2)_{\mathrm{R}} \otimes U(1)_{B-L} \\ &\xrightarrow{\chi_{\mathrm{R}}} SU(3)_C \otimes SU(2)_{\mathrm{L}} \otimes U(1)_Y \\ &\xrightarrow{\chi_{\mathrm{L}}} SU(3)_C \otimes U(1)_{\mathrm{em}} . \end{aligned} \tag{15}$$

The Higgs fields are given in Table 2. The component $[\mathbf{1},\mathbf{1},\mathbf{1}]$ of $\{\mathbf{210}\}$ is $\mathcal{D}$-odd.

If we use only the fields of $\{\mathbf{210}\}$ in order to produce the first two steps in (15), then the potential analogous to (12) has to be modified to

$$\begin{aligned} \mathcal{L} = &\mu^2(\mathbf{16}^* \times \mathbf{16}) + \lambda_\chi(\mathbf{16}^* \times \mathbf{16})^2 + m_M^2(\mathbf{210})^2 \\ &+ \eta_M(\mathbf{210})^3 + \lambda_M(\mathbf{210})^4 + M_M(\mathbf{210})(\mathbf{16}^* \times \mathbf{16}) \\ &+ \varepsilon_M(\mathbf{210})^2 \times (\mathbf{16}^* \times \mathbf{16}) + \kappa[(\mathbf{16}^*)^4 + (\mathbf{16})^4] . \end{aligned} \tag{16}$$

Then, following [15], the masses of the Higgs doublets obtained from (16) are given by

$$\begin{aligned} m_{\mathrm{R}}^2 &= \mu^2 + M_M s_M + \varepsilon_M s_M^2 , \\ m_{\mathrm{L}}^2 &= \mu^2 - M_M s_M + \varepsilon_M s_M^2 . \end{aligned} \tag{17}$$

As we have $\langle\chi_{\mathrm{L}}\rangle = \upsilon_{\mathrm{L}} \sim m_{\mathrm{L}}$ and $\langle\chi_{\mathrm{R}}\rangle = \upsilon_{\mathrm{R}} \sim m_{\mathrm{R}}$, we find, after tuning of the model parameters, that $\upsilon_{\mathrm{L}} \sim 100$ GeV and $\upsilon_{\mathrm{R}} \sim$ GUT. In fact, we have the relation

$$\upsilon_{\mathrm{R}}^2 - \upsilon_{\mathrm{L}}^2 = \frac{M_M s_M}{2\kappa} . \tag{18}$$

From this equation we see that, due to the breaking of $\mathcal{D}$-parity in the GUT scale by the field $\mathcal{D}$-odd $[\mathbf{1},\mathbf{1},\mathbf{1}] \subset \{\mathbf{210}\}$, the breaking of the L–R symmetry is also induced close to the GUT scale. This is a prediction different from the model given by the breaking chain (11), in which it is possible that the breaking of L–R symmetry occurs close to the TeV scale. Other differences are also possible from the renormalization groups equations (RGE). This analysis for the breaking chain (15) was done in [13] at one and two loops on the gauge couplings. Some stages of the breaking chain (11) have been analyzed in different papers [16, 17] in the same context of the RGE. In fact, the use of the G_{SM} singlets in the $\{\mathbf{16}\}$, $\{\mathbf{126}\}$ and $\{\mathbf{144}\}$ representations would lead to the result that the unification of the G_{SM} coupling constants is inconsistent with the low-energy data on these couplings. On the other side, the use of $\{\mathbf{54}\}$ in the breaking of $SO(10) \xrightarrow{S_M} G_{\mathrm{PS}} \otimes \mathcal{D}$ seems to be consistent with the experimental bound on the proton lifetime only at a marginal level [16, 17]. We conclude that a more complete analysis of the RGE would be necessary for the breaking chains along the lines of [18].

4 An L–R model based in $SU(7)$ with mirrorfermions

In the L–R model with mirror fermions the particle content is described in Table 3 for the two first families with its quantum numbers under $SU(3)_C \otimes SU(2)_{\mathrm{L}} \otimes SU(2)_{\mathrm{R}} \otimes U(1)_Y$.

Some points should be observed. First, as $SU(3)_C \otimes SU(2)_{\mathrm{L}} \otimes U(1)_Y$ is a maximal subgroup of $SU(5)$, then we have $SU(3)_C \otimes SU(2)_{\mathrm{L}} \otimes SU(2)_{\mathrm{R}} \otimes U(1)_Y \subset SU(5) \otimes SU(2)_{\mathrm{R}} \subset SU(7)$. In fact [19] $SU(5) \otimes SU(2) \otimes U(1)_X$ is a maximal subgroup of $SU(7)$ and we can assume $SU(2)$ to have the right chirality $SU(2)_{\mathrm{R}}$.

A second point is that the mass terms of leptons $\overline{l_{e_{\mathrm{L}}}}\chi_{\mathrm{L}} e_{\mathrm{R}}$ require Higgs representations $\chi_{\mathrm{L}} \sim (\mathbf{1},\mathbf{2},\mathbf{1},1)$. Similarly the mass terms of the mirror partners $\overline{L_{E_{\mathrm{R}}}}\chi_{\mathrm{R}} E_{\mathrm{L}}$, require $\chi_{\mathrm{R}} \sim (\mathbf{1},\mathbf{1},\mathbf{2},1)$. Mixing terms of the type $\overline{e_{\mathrm{R}}} S_D E_{\mathrm{L}}$, $\overline{\nu_{\mathrm{R}}} S_D N_{EL}$ need $S_D \sim (\mathbf{1},\mathbf{1},\mathbf{1},0)$. Mass terms of the Majorana type $\overline{l_{e_{\mathrm{L}}}}\widetilde{\chi_{\mathrm{L}}} N_{EL}^C$ need $\widetilde{\chi_{\mathrm{L}}} \sim (\mathbf{1},\mathbf{2},\mathbf{1},-1)$, and $\overline{L_{E_{\mathrm{R}}}}\widetilde{\chi_{\mathrm{R}}}\nu_{e_{\mathrm{R}}}^C$ need $\widetilde{\chi_{\mathrm{R}}} \sim (\mathbf{1},\mathbf{1},\mathbf{2},-1)$ in order to give mass to neutrinos. The $\overline{N_{EL}^C} S_M N_{EL}$ and $\overline{\nu_{e_{\mathrm{R}}}^C} S_M \nu_{e_{\mathrm{R}}}$ terms are possible with $S_M \sim (\mathbf{1},\mathbf{1},\mathbf{1},0)$. Now, let us search for the

Table 3.

Ordinary fermions	Mirror fermions
$l_L = \begin{pmatrix} \nu_e \\ e \end{pmatrix}_L, \begin{pmatrix} \nu_\mu \\ \mu \end{pmatrix}_L \sim (\mathbf{1},\mathbf{2},\mathbf{1},-1)$	$L_R = \begin{pmatrix} N_E \\ E \end{pmatrix}_R, \begin{pmatrix} N_M \\ M \end{pmatrix}_R \sim (\mathbf{1},\mathbf{1},\mathbf{2},-1)$
$e_R, \mu_R \sim (\mathbf{1},\mathbf{1},\mathbf{1},-2)$	$E_L, M_L \sim (\mathbf{1},\mathbf{1},\mathbf{1},-2)$
$\nu_{eR}, \nu_{\mu R} \sim (\mathbf{1},\mathbf{1},\mathbf{1},0)$	$N_{EL}, N_{ML} \sim (\mathbf{1},\mathbf{1},\mathbf{1},0)$
$\begin{pmatrix} u \\ d \end{pmatrix}_L, \begin{pmatrix} c \\ s \end{pmatrix}_L \sim (\mathbf{3},\mathbf{2},\mathbf{1},1/3)$	$\begin{pmatrix} U \\ D \end{pmatrix}_R, \begin{pmatrix} C \\ S \end{pmatrix}_R \sim (\mathbf{3},\mathbf{1},\mathbf{2},1/3)$
$u_R, c_R \sim (\mathbf{3},\mathbf{1},\mathbf{1},4/3)$	$U_L, C_L \sim (\mathbf{3},\mathbf{1},\mathbf{1},4/3)$
$d_R, s_R \sim (\mathbf{3},\mathbf{1},\mathbf{1},-2/3)$	$D_L, S_L \sim (\mathbf{3},\mathbf{1},\mathbf{1},-2/3)$

representations of $\chi_{\mathrm{L,R}}$, S_D and S_M in the $SU(7)$ context [20, 21]. The fermionic multiplet is in the anomaly free combination[1] [22] $\{\mathbf{1}\}\oplus\{\mathbf{7}\}\oplus\{\mathbf{21}\}\oplus\{\mathbf{35}\}$, corresponding to the spinor representation **64** of $SO(14)$ in which $SU(7)$ is embedded. In the previous multiplets, $\{\mathbf{21}\}$ is a 2-fold, $\{\mathbf{35}\}$ is a 4-fold and $\{\mathbf{7}\}$ is a 6-fold of totally antisymmetric tensors.

Let us note that **64** can contain two families of ordinary fermions with its respective mirror partners, for example the electron and muon families as is shown in Table 3. The other families can be incorporated into the other **64** spinorial representation. The branching rules for each component of the spinorial representation, under its subgroup $SU(5)\otimes SU(2)_{\mathrm{R}}$, are [23–25]

$$\begin{aligned}\{\mathbf{35}\} &= [\mathbf{10}^*,\mathbf{1}]\oplus[\mathbf{10},\mathbf{2}]\oplus[\mathbf{5},\mathbf{1}]\,,\\ \{\mathbf{21}\} &= [\mathbf{10}^*,\mathbf{1}]\oplus[\mathbf{5}^*,\mathbf{2}]\oplus[\mathbf{1},\mathbf{1}]\,,\\ \{\mathbf{7}\} &= [\mathbf{5},\mathbf{1}]\oplus[\mathbf{1},\mathbf{2}]\,,\end{aligned}\tag{19}$$

and under $SU(3)_C\otimes SU(2)_{\mathrm{L}}\otimes SU(2)_{\mathrm{R}}\otimes U(1)_Y$ they are

$$\begin{aligned}\{\mathbf{35}\} = {} & \underbrace{(\mathbf{1},\mathbf{1},\mathbf{1},-2)}_{e_{\mathrm{R}}}\oplus\underbrace{(\mathbf{3},\mathbf{1},\mathbf{1},4/3)}_{u_{\mathrm{R}}}\oplus\underbrace{(\mathbf{3},\mathbf{2},\mathbf{1},1/3)}_{\begin{pmatrix}c\\ s\end{pmatrix}_{\mathrm{L}}}\\ & \oplus\underbrace{(\mathbf{1},\mathbf{1},\mathbf{1},-2)}_{E_{\mathrm{L}}}\oplus\underbrace{(\mathbf{1},\mathbf{1},\mathbf{1},-2)}_{M_{\mathrm{L}}}\oplus\underbrace{(\mathbf{3},\mathbf{1},\mathbf{1},4/3)}_{U_{\mathrm{L}}}\\ & \oplus\underbrace{(\mathbf{3},\mathbf{1},\mathbf{1},4/3)}_{C_{\mathrm{L}}}\oplus\underbrace{(\mathbf{3},\mathbf{1},\mathbf{1},-2/3)}_{s_{\mathrm{R}}}\oplus\underbrace{(\mathbf{1},\mathbf{2},\mathbf{1},-1)}_{\begin{pmatrix}\nu_e\\ e\end{pmatrix}_{\mathrm{L}}}\\ & \oplus\underbrace{(\mathbf{3},\mathbf{1},\mathbf{2},1/3)}_{\begin{pmatrix}U\\ D\end{pmatrix}_{\mathrm{R}}}\oplus\underbrace{(\mathbf{3},\mathbf{1},\mathbf{2},1/3)}_{\begin{pmatrix}C\\ S\end{pmatrix}_{\mathrm{R}}}\,,\end{aligned}\tag{20}$$

$$\begin{aligned}\{\mathbf{21}\} = {} & \underbrace{(\mathbf{1},\mathbf{1},\mathbf{1},-2)}_{\mu_{\mathrm{R}}}\oplus\underbrace{(\mathbf{3},\mathbf{1},\mathbf{1},4/3)}_{c_{\mathrm{R}}}\oplus\underbrace{(\mathbf{3},\mathbf{2},\mathbf{1},1/3)}_{\begin{pmatrix}u\\ d\end{pmatrix}_{\mathrm{L}}}\\ & \oplus\underbrace{(\mathbf{1},\mathbf{1},\mathbf{2},-1)}_{\begin{pmatrix}N_E\\ E\end{pmatrix}_{\mathrm{R}}}\oplus\underbrace{(\mathbf{1},\mathbf{1},\mathbf{2},-1)}_{\begin{pmatrix}N_M\\ M\end{pmatrix}_{\mathrm{R}}}\oplus\underbrace{(\mathbf{3},\mathbf{1},\mathbf{1},-2/3)}_{D_{\mathrm{L}}}\\ & \oplus\underbrace{(\mathbf{3},\mathbf{1},\mathbf{1},-2/3)}_{S_{\mathrm{L}}}\oplus\underbrace{(\mathbf{1},\mathbf{1},\mathbf{1},0)}_{N_{M_{\mathrm{L}}}}\,,\end{aligned}\tag{21}$$

$$\begin{aligned}\{\mathbf{7}\} = {} & \underbrace{(\mathbf{1},\mathbf{2},\mathbf{1},-1)}_{\begin{pmatrix}\nu_\mu\\ \mu\end{pmatrix}_{\mathrm{L}}}\oplus\underbrace{(\mathbf{3},\mathbf{1},\mathbf{1},-2/3)}_{d_{\mathrm{R}}}\\ & \oplus\underbrace{(\mathbf{1},\mathbf{1},\mathbf{1},0)}_{N_{EL}}\oplus\underbrace{(\mathbf{1},\mathbf{1},\mathbf{1},0)}_{N_{M_{\mathrm{L}}}}\,.\end{aligned}\tag{22}$$

$$\{\mathbf{1}\} = \underbrace{(\mathbf{1},\mathbf{1},\mathbf{1},0)}_{\nu_{e_{\mathrm{R}}}}\,.\tag{23}$$

From the product $\{\mathbf{63}\}\otimes\{\mathbf{63}\}=\{\mathbf{1}\}\oplus\{\mathbf{63}\}\oplus\ldots$, we obtain the Higgs representations producing the mass terms for the fermions in the spinorial multiplet $\{\mathbf{63}\}=\{\mathbf{7}\}\oplus\{\mathbf{21}\}\oplus\{\mathbf{35}\}$ of $SU(7)$. With the help of the branching rules (15)–(17), we take

$$\chi_{\mathrm{L}}\sim\{\mathbf{7}^*\}\supset(\mathbf{1},\mathbf{2},\mathbf{1},1)\,,\quad \chi_{\mathrm{R}}\sim\{\mathbf{21}^*\}\supset(\mathbf{1},\mathbf{1},\mathbf{2},1)\,,\tag{24}$$

$$S_D\sim\{\mathbf{21}\}\supset(\mathbf{1},\mathbf{1},\mathbf{1},0)\,,\ S_M\sim\{\mathbf{1}\sim(\mathbf{1},\mathbf{1},\mathbf{1},0)\,.\tag{25}$$

Finally we can have the following breaking chain with two singlets and two doublets of Higgs representations:

$$\begin{aligned}SU(7) &\xrightarrow{S_M} SU(5)\otimes SU(2)_{\mathrm{R}}\otimes\mathcal{D}\\ &\xrightarrow{S_D} G_{\mathrm{SM}}\otimes SU(2)_{\mathrm{R}}\\ &\xrightarrow{\chi_{\mathrm{R}}} SU(3)_C\otimes SU(2)_{\mathrm{L}}\otimes U(1)_Y\\ &\xrightarrow{\chi_{\mathrm{L}}} SU(3)_C\otimes U(1)_{\mathrm{em}}\,.\end{aligned}\tag{26}$$

The component of $\phi^{\alpha\beta}=\{\mathbf{21}\}$ that breaks $\mathcal{D}$-parity is given by $S_D=\phi^{67}$, which is odd under $\mathcal{D}$-parity [26], and the S_M field, being an $SU(7)$ singlet, preserves $\mathcal{D}$-parity.

We can write an $SU(7)$ invariant Higgs potential similar to the one given in (12) with the obvious changes $\mathbf{144}\longrightarrow\{\mathbf{7}^*\}$, $\mathbf{45}\longrightarrow\{\mathbf{21}\}$, and $\mathbf{54}\longrightarrow\{\mathbf{1}\}$.

5 Conclusions

In conclusion we have shown that parity can be spontaneously broken by a simple Higgs sector with two doublets and two singlets. The grand unified sector that contains this possibility is more restricted than other scenarios. One of the new proposed singlets can have a breaking scale not very far from the Fermi scale. A similar conclusion was recently found in a different approach for two doublets models [27]. In the case of the model with mirror fermions a significant contribution is possible to the magnetic moment of electrons and muons [14] due to couplings with the mirror fermions of the type $f(\overline{l_{\mathrm{L}}}\chi_{\mathrm{L}}e_{\mathrm{R}}+\overline{l_{\mathrm{R}}}\chi_{\mathrm{R}}E_{\mathrm{L}})+f'\overline{e_{\mathrm{R}}}E_{\mathrm{L}}S_D$. These terms can give an important contribution to the muon anomaly and will be connected to the breaking of the Weinberg symmetry[2].

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[1] We are using { } for the $SU(7)$ components too.

[2] Work in this direction was done in the E_6 model; see [28].

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