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The Linear BESS model at the LHC

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Received 6 May 2018; received in revised form 16 July 2018; accepted 18 July 2018

Editor: Tommy Ohlsson

Abstract

In this work we consider the Linear BESS model at the LHC. This model can be seen as an adequated benchmark for exploring the phenomenological consequences of a composite Higgs sector since its particle content is the one we would expect in a realistic low energy description of modern (Technicolor inspired) dynamical electroweak symmetry breaking scenarios. Additionally, the model exhibits the property of decoupling, producing a good ultraviolet behavior. We focus on the limits on the masses of the new heavy vector particles imposed by direct resonance searches, recent measurements of the decay of the Higgs boson into two photons and the electroweak precision tests. We found that the model is capable to accommodate the existing experimental constrains provided that the spin-1 resonances are heavier than 3.4 TeV.

1. Introduction

Despite its enormous experimental success, the Standard Model (SM) offers some aspects which are not completely satisfactory from the theoretical point of view. One of them is that not all the interactions present in the model have their origin in the gauge principle. Indeed, the scalar potential and the Yukawa interactions, which are central parts of the model, are not dictated by a local symmetry. In fact, it may be argued that this is the origin of crucial problems of the SM like the Naturalness Problem. A very elegant solution to this criticism is to assume that the Higgs sector and the Yukawa interactions have their origin in a strongly interacting

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https://doi.org/10.1016/j.nuclphysb.2018.07.014
0550-3213/

Please cite this article in press as: J. Urbina, A.R. Zerwekh, The Linear BESS model at the LHC, Nucl. Phys. B (2018), https://doi.org/10.1016/j.nuclphysb.2018.07.014

gauge theory. A paradigmatic example of this framework is (extended) Technicolor, specially its modern incarnation: Walking Technicolor. When working with strongly interacting theories, it is usually convenient to consider effective approaches which include only the (composite) degrees of freedom which are relevant in the low energy limit. A very well known effective approach to the Dynamical Electroweak Symmetry Breaking paradigm is the BESS model [1,2]. Originally, the BESS model (as the old Technicolor idea) was Higgsless. However, the Higgs boson was discovered and, consequently, the original version is ruled out. Fortunately, toward the end of the 1990's, a version of the BESS model which included scalar fields was formulated. This is the so called Linear BESS (LBESS) model [3,4]. This model is of special interest because its particle content (two isotriplet spin-1 resonances, a Higgs-like scalar and two heavy scalars) is the one we would expect in a realistic low energy description of a new strong sector responsible for the electroweak symmetry breaking [5,6]. Additionally, the LBESS model is renormalizable and it possesses the property of decoupling [3,4], *i.e.*, for a large New Physics scale the SM is exactly recovered. This is a very important feature, given the current lack of unequivocal signals beyond SM at the LHC, which indicates that the scale of New Physics (if any) may be larger than originally expected. Indeed, the LBESS model may provide a useful interpolation in the phenomenological studies of two kinds of Dynamical Electroweak Symmetry Breaking scenarios: one characterized by a low scale of a few TeV (like Walking Technicolor) and a second one where the Higgs boson appears as a pseudo-Nambu-Goldstone boson (the so called Composite Higgs models) and the typical New Physics scale is of the order of 10 TeV. In summary, the LBESS model helps us to study the phenomenological implications of scenarios which elegantly explain the dynamics underlying the electroweak symmetry breaking process and solve the Naturalness problem, and simultaneously promises to be consistent with the current experimental situation.

In this work, we study the phenomenology of this model at the LHC. We focus on four kinds of measurements in order to constrain the parameter space of the model: the Higgs decay into a pair of photons, resonance searches in the dijet and the dilepton spectra, and the precision electroweak tests.

The paper is organized as follows. In section 2 we briefly recall the main features of the LBESS model. In section 3, we describe our simulations and results, while in section 4 we state our conclusions. For the sake of completeness we add a longer description of the model in an appendix.

2. Recalling the Linear BESS model

The basic point of view behind the BESS model attributes the origin of the electroweak scale to a new strongly interacting sector in analogy to the dynamical origin of Λ_{QCD} in QCD. The hypothetical new strong interaction is supposed to be confining and at low energy it manifests itself through composite states. It is expected that the lightest composite particles be scalars and vector resonances. In this section, we provide a general description of the main features of the model. More details can be found in the Appendix and in the original literature [1–4].

Following the Hidden Local Symmetry (HLS) formalism, the composite vectors can be introduced as gauge fields of effective gauge groups. Consequently, in the LBESS model, we start with an extended gauge symmetry given by $SU(2)_L \otimes U(1) \otimes SU(2)'_L \otimes SU(2)'_R$. The symmetry is broken down to $U(1)_{\text{em}}$ in two steps as shown in the following scheme

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$$\begin{array}{c} SU(2)_L \otimes U(1) \otimes SU(2)'_L \otimes SU(2)'_R \\ \downarrow u \\ SU(2)_{\text{weak}} \otimes U(1)_Y \\ \downarrow v \\ U(1)_{\text{em}} \end{array}$$

where u is the scale characterizing the breaking of the HLS and v is the usual electroweak scale. All the symmetry breaking processes are assumed to be produced by the vacuum expectation values of (composite) scalar fields $\langle \rho_U \rangle = v$ and $\langle \rho_L \rangle = \langle \rho_R \rangle = u$. The breaking down of the symmetries produces non-diagonal mass matrices in the gauge and scalar sectors. The physical spectrum (composed by the mass eigenstates) consists on the following fields:

1. *Two heavy vector triplet*: (V_L^+, V_L^0, V_L^-) and (V_R^+, V_R^0, V_R^-). These vector bosons are mainly the gauge bosons of $SU(2)'_L \otimes SU(2)'_R$ with a small mixing with the gauge field of $SU(2)_L \otimes U(1)$. Naturally, they have masses of the order of $g_2 u$ where g_2 is the coupling constant associated with the groups $SU(2)'_L$ and $SU(2)'_R$. Because the standard fermions are assumed to be charged only under $SU(2)_L$ (left-handed) and $U(1)$ (left-handed and right-handed), the heavy vectors V_L^+, V_L^0, V_L^- and V_R^0 couple to the standard fermions with coupling constants proportional to the (small) mixing angles. Notice that V_R^+ and V_R^- do not couple to the standard fermions.
2. *The standard electroweak gauge bosons*: $W^\pm, Z, A$.
3. *Two heavy scalars: H_L, H_R*. These scalars are supposed to have masses of the order of the u scale. They correspond mainly to the original ρ_L and ρ_R fields, respectively, with and small mixing with ρ_U. Originally, the standard fermions can form Yukawa terms only with ρ_U due to the quantum numbers assigned to the fermions. This means that H_L and H_R coupling to the standard fermions is proportional to small mixing angles.
4. *The standard-like Higgs boson: H* which correspond mainly to the original ρ_U field.

The model has six free parameters, namely: the masses of the heavy vectors (M_{V_L} and M_{V_R}), the masses of the heavy scalar (M_{H_L} and M_{H_R}), the scale of the HLS breakdown (u) and a parameter governing a quartic interaction term between scalars (f).

In what follows, we will assume that u, M_{H_L} and M_{H_R} are of the order of 3 TeV while the masses of the heavy vectors will be taken in the range of 2 to 4 TeV. The assumption of very massive scalars beside a light Higgs-like boson is well justified in this particular model (see equation (11) in the Appendix) and has also been found to be self-consistent in a similar effective model previously studied by our group [6].

3. Results

3.1. $H \to \gamma\gamma$

The first process we consider is the Higgs boson decay into two photons. This is a 1-loop process which includes the contribution of the new charged states: in our case, the new charged vector bosons $V_L^\pm$ and $V_R^\pm$. For heavy vector bosons with moderate coupling to the Higgs boson, it is expected that this process does not deviate significantly from the SM [6,7]. This is exactly our case: the new vector bosons, as described above, are considered in the 2–4 TeV mass range. On the other hand, the coupling between the Higgs and $V_R^\pm$ originates from the mixing of the different scalar fields of the model. This mixing (and thus the referred coupling) is controlled

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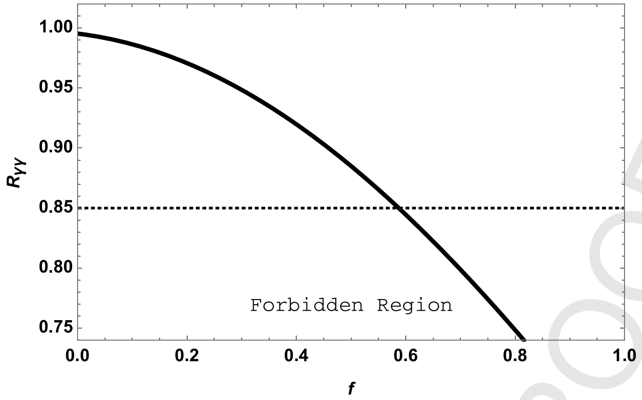

Fig. 1. Predicted values of $R_{\gamma\gamma}$ (continuous line), as a function of the f parameter, compared to the lower limit of the experimental value at 1σ (dashed line). The region above the dashed line is allowed.

by the f parameter of the scalar potential, which has to be positive. In fact, we found that this parameter is the only one sensible to the current data for this process. It is useful to define the ratio

$$R_{\gamma\gamma} = \frac{\sigma\,(pp \to H)\,\Gamma\,(H \to \gamma\gamma)}{\sigma\,(pp \to H)_{\mathrm{SM}}\,\Gamma\,(H \to \gamma\gamma)_{\mathrm{SM}}} = \frac{\Gamma\,(H \to \gamma\gamma)}{\Gamma\,(H \to \gamma\gamma)_{\mathrm{SM}}} \tag{1}$$

in order to quantify the departure of the model from the SM's predictions. The second equality in (1) is due to the fact that the Higgs production mechanism is not modified by the LBESS model. The most recent measurement of $R_{\gamma\gamma}$ has been done by ATLAS and results to be $R_{\gamma\gamma} = 0.99 \pm 0.14$ [8]. In Fig. 1 we show the values of $R_{\gamma\gamma}$ predicted by the LBESS model as a function of the f parameter (continuous line). For comparison purpose, we also include the lowest limit (at 1σ) of the experimental value. We can see that the model is in good agreement with experiment for $f \in [0, 0.6]$.

3.2. Searching for resonances

We also consider the direct searches of resonances in the dijet and the dilepton channels. In the kinematic setup we have adopted (with very heavy non-standard scalars) only the spin-1 particle $V_L^{\pm}$, V_L^0 (assumed to be degenerated) and V_R^0 can be produced in the s-channel by quark–anti-quark annihilation. At this point, we recall that the fields $V_R^{\pm}$ do not couple to the standard fermions. On the other hand, in the construction we are considering, V_R cannot be heavier than V_L (see equation (16)). Consequently, the model predicts that two resonances should appear in the dijet and the dilepton spectra with the lighter one corresponding to V_R^0.

3.2.1. Methodology

We implemented the model in CalcHEP [9] using the LanHEP package [10,11] and generated events for dijet and dilepton production at the 13 TeV LHC, considering only the contribution of the new particles, without background, for several values of M_{V_L} and M_{V_R}. As an example, in Fig. 2 we show two dijet spectra obtained in our simulations. In order to put constrains on the model parameter space, we count the events associated to the peaks, we compute a cross section for each resonance and we compare it with the experimental upper limits for dijet or dilepton

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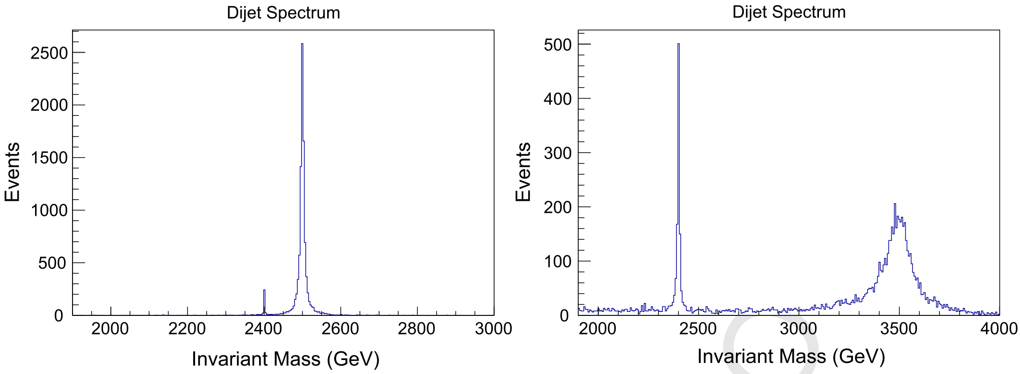

Fig. 2. Examples of resonances in the dijet invariant mass spectrum. In these simulations has been taken into account only the contribution of the non-standard sector of the LBESS model without background.

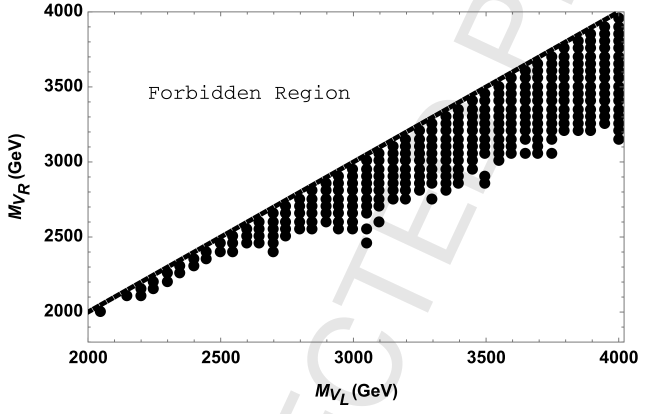

Fig. 3. Constrains to the (M_{V_L}, M_{V_R}) space by direct resonance searches in the dijets spectrum at the 13 TeV LHC. The region filled with black circles is the allowed one.

resonance cross-sections and we only accept a pair of M_{V_L} and M_{V_R} values if the cross-section associated to each resonance is smaller than the experimental limit.

3.2.2. Dijets

In the case of dijets, we use the limits provided by the ATLAS Collaboration based on data taken at $\sqrt{s} = 13$ TeV [12]. Our results are shown in Fig. 3. The region filled with black circles is the accepted zone, that is, the set of points (M_{V_L}, M_{V_R}) which produce both resonances with a cross sections below the experimental limit. Notice that in the quasi-degenerated case, resonances with masses as light as 2000 or 2500 GeV are allowed. This is in agreement with previous studies on vector resonances [6,7,13].

3.2.3. Dileptons

More stringent constrains are obtained in the dilepton channel. In this case, we use the limits provided by the ATLAS Collaboration at $\sqrt{s} = 13$ TeV and 36.1 fb^{-1} [14]. Again, the region filled by black circles is the accepted zone (Fig. 4). In this case, only resonances heavier than 3.4 TeV are allowed. This improvement on the constrains is mainly due to the fact that the

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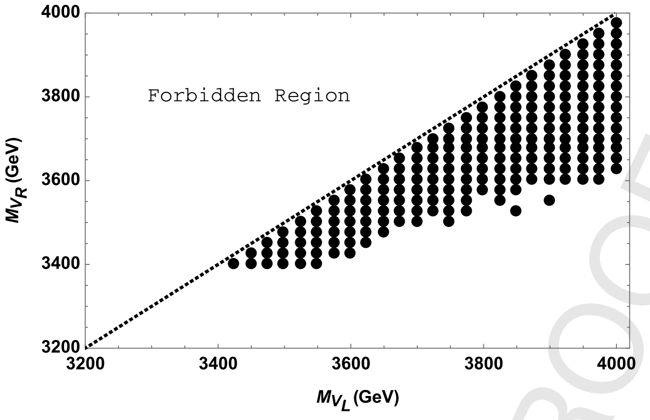

Fig. 4. Constrains to the (M_{V_L}, M_{V_R}) space by direct resonance searches in the dilepton spectrum at the 13 TeV LHC. The region filled with black circles is the allowed one.

dilepton production is a cleaner channel than dijet production at a hadron collider and the higher luminosity of the dilepton set of data.

3.3. Precision tests

An indirect way to constrain extensions of the SM is to consider the contribution the new Physics provides to the electroweak radiative corrections. These effects are, in general, parametrized by the well known Peskin–Takeuchi parameters: S, T and U. An equivalent set of parameters, named ϵ_1, ϵ_2 and ϵ_3, is also used in the literature [15] and is related to the former one by the following expressions:

$$\begin{aligned} \epsilon_1 &= \alpha T \\ \epsilon_2 &= -\frac{\alpha}{4s_Z^2} U \\ \epsilon_3 &= \frac{\alpha}{4s_Z^2} S \end{aligned} \tag{2}$$

where α is the electromagnetic coupling constant at the scale of M_Z and s_Z^2 is the $\sin^2\theta_W$ in the $\overline{MS}$ scheme at the same scale. In Ref. [4], the following tree level expressions are provided for the ϵ_i parameters in the context of the LBESS model:

$$\begin{aligned} \epsilon_1 &= -r s_\varphi^2 \frac{c_\theta^4 + s_\theta^4}{c_\theta^4} \\ \epsilon_2 &= -r s_\varphi^2 \\ \epsilon_3 &= -\frac{r s_\varphi^2}{c_\theta^2} \end{aligned} \tag{3}$$

where $c_\theta = \cos\theta_W$, $s_\theta = \sin\theta_W$, while r and s_φ can be expressed in terms of M_{V_L} and M_{V_R} as shown in the Appendix.

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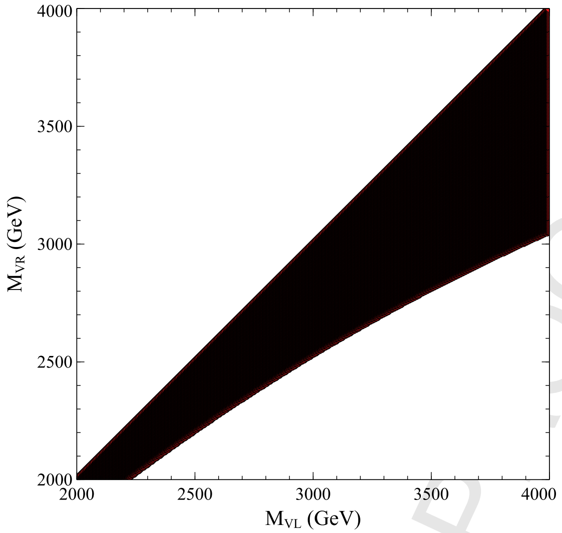

Fig. 5. Constrains to the (M_{V_L}, M_{V_R}) space by the electroweak precision tests. The black region is the allowed one.

We use the values $S = 0.05 \pm 0.10$, $T = 0.08 \pm 0.12$ and $U = 0.02 \pm 0.10$ found in [16], and the expressions above to select the combination of M_{V_L} and M_{V_R} which reproduce simultaneously the experimental values of ϵ_1, ϵ_2 and ϵ_3 (using equations (3)) within 1σ. The result is shown in Fig. 5. The black region is the zone allowed by the precision variables. As we can see, the restrictions are not as stringent as the ones obtained using dilepton data.

4. Conclusions

We have studied the LBESS model in the context of recent data obtained at the 13 TeV LHC in the regime where the non-standard scalars are heavy with masses of the order of 3 TeV. We found that the value the f parameter of the scalar potential has to belong to the interval $[0, 0.6]$. Additionally, we have found that the model is consistent with current experimental data provided that the spin-1 resonances are heavier than 3.4 TeV. This limit is higher than the ones obtained in previous studies of models with vector resonances. The main restrictions come from recent searches of resonances in the dilepton spectrum.

Acknowledgements

The authors want to thank Bastián Díaz and Felipe Rojas for useful discussions.

This work was supported in part by CONICYT (Chile) grants PIA/ACT-1406 and PIA/Basal FB0821, and by FONDECYT (Chile) grant 1160423.

Appendix

The LBESS model is based on the global symmetry $SU(2)_L \otimes SU(2)_R \otimes SU(2)'_L \otimes SU(2)'_R$ from which the subgroup $SU(2)_L \otimes U(1) \otimes SU(2)'_L \otimes SU(2)'_R$ is made local. The matter content consists of three scalars (Φ_U, Φ_L and Φ_R) which we are supposed to be composite and the

Table 1
Representations of the scalar fields under the full global group.

Fields	$SU(2)_L$	$SU(2)_R$	$SU(2)'_L$	$SU(2)_R$
Φ_{U}	**2**	$\bar{\mathbf{2}}$	1	1
Φ_{L}	**2**	1	**2**	1
Φ_{R}	1	**2**	1	**2**

Table 2
Representations of the fermion fields under the full global group.

Fermion	$SU(2)_L$	$SU(2)_R$	$SU(2)'_L$	$SU(2)'_R$
Ψ_{iL}	**2**	1	1	1
Ψ_{iR}	1	**2**	1	1

Table 3
Representations of the fermion fields under the full gauge group. Here $Y' = B - L$ with B and L being the baryon and lepton numbers.

Fermion	$SU(2)_L$	$U(1)$	$SU(2)'_L$	$SU(2)'_R$
Ψ_{iL}	**2**	Y'	1	1
Ψ_{iR}	1	Y'	1	1

standard fermions (plus a right-handed neutrino). These fields transform under the global symmetry following the representation assignment showed in Tables 1 (scalars) and 2 (fermions). The representation assignment for fermions under the local symmetry is shown in Table 3.

The full Lagrangian of model can be written down as:

$$
\begin{aligned}
\mathcal{L} = & \frac{1}{2}\mathrm{Tr}[F_{\mu\nu}(\tilde{W})F^{\mu\nu}(\tilde{W})] + \frac{1}{2}\mathrm{Tr}[F_{\mu\nu}(\tilde{B})F^{\mu\nu}(\tilde{B})] \\
& + \frac{1}{2}\mathrm{Tr}[F_{\mu\nu}(A_L)F^{\mu\nu}(A_L)] + \frac{1}{2}\mathrm{Tr}[F_{\mu\nu}(A_R)F^{\mu\nu}(A_R)] \\
& + \frac{1}{4}\left[\mathrm{Tr}[(D_\mu\Phi_{\mathrm{U}})^\dagger(D^\mu\Phi_{\mathrm{U}})] + \mathrm{Tr}[(D_\mu\Phi_{\mathrm{L}})^\dagger(D^\mu\Phi_{\mathrm{L}})] + \mathrm{Tr}[(D_\mu\Phi_{\mathrm{R}})^\dagger(D^\mu\Phi_{\mathrm{R}})]\right] \\
& - \mu^2(\mathrm{Tr}[\Phi_{\mathrm{L}}^\dagger\Phi_{\mathrm{L}}] + \mathrm{Tr}[\Phi_{\mathrm{R}}^\dagger\Phi_{\mathrm{R}}]) - \frac{\lambda}{4}([\mathrm{Tr}(\Phi_{\mathrm{L}}^\dagger\Phi_{\mathrm{L}})]^2 \\
& + [\mathrm{Tr}(\Phi_{\mathrm{R}}^\dagger\Phi_{\mathrm{R}})]^2) - m^2\mathrm{Tr}(\Phi_{\mathrm{U}}^\dagger\Phi_{\mathrm{U}}) - \frac{h}{4}[\mathrm{Tr}(\Phi_{\mathrm{U}}^\dagger\Phi_{\mathrm{U}})]^2 \\
& - \frac{f_3}{2}\mathrm{Tr}(\Phi_{\mathrm{L}}^\dagger\Phi_{\mathrm{L}})\mathrm{Tr}(\Phi_{\mathrm{R}}^\dagger\Phi_{\mathrm{R}}) - \frac{f}{2}[\mathrm{Tr}(\Phi_{\mathrm{L}}^\dagger\Phi_{\mathrm{L}})\mathrm{Tr}(\Phi_{\mathrm{U}}^\dagger\Phi_{\mathrm{U}}) + \mathrm{Tr}(\Phi_{\mathrm{R}}^\dagger\Phi_{\mathrm{R}})\mathrm{Tr}(\Phi_{\mathrm{U}}^\dagger\Phi_{\mathrm{U}})] \\
& + i\bar{\Psi}_{iL}\gamma^\mu D_\mu\Psi_{iL} + i\bar{\Psi}_{iR}\gamma^\mu D_\mu\Psi_{iR} + \bar{\Psi}_{iL}\Phi_{\mathrm{U}}\Upsilon\Psi_{jR} + \mathrm{h.c.}
\end{aligned}
\tag{4}
$$

where $\tilde{W}_\mu$, $\tilde{B}_\mu$, $A_{L\mu}$ and $A_{R\mu}$ are the gauge bosons associated to $SU(2)_L$, $U(1)$, $SU(2)'_L$ and $SU(2)'_R$, respectively and Υ is a matrix containing the Yukawa coupling constants. On the other hand, the strength tensors are defined as usual:

$$
\begin{aligned}
F_{\mu\nu}(\tilde{W}) &= \partial_\mu\tilde{W}_\nu - \partial_\nu\tilde{W}_\mu + g_0[\tilde{W}_\mu, \tilde{W}_\nu] \\
F_{\mu\nu}(\tilde{B}) &= \partial_\mu\tilde{B}_\nu - \partial_\nu\tilde{B}_\mu \\
F_{\mu\nu}(A_L) &= \partial_\mu A_{L\nu} - \partial_\nu A_{L\mu} + g_2[A_{L\mu}, A_{L\nu}]
\end{aligned}
$$

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$$F_{\mu\nu}(A_R) = \partial_\mu A_{R\nu} - \partial_\nu A_{R\mu} + g_2[A_{R\mu}, A_{R\nu}] \tag{5}$$

The covariant derivatives in the kinetic terms for the scalars and fermions in equation (4), are given by:

$$\begin{aligned}
D_\mu \Phi_\mathrm{L} &= \partial_\mu \Phi_\mathrm{L} + i g_0 \frac{\tau^a}{2} \tilde{W}^a_\mu \Phi_\mathrm{L} - i g_2 \Phi_\mathrm{L} \frac{\tau^a}{2} A^a_{L\mu} \\
D_\mu \Phi_\mathrm{R} &= \partial_\mu \Phi_\mathrm{R} + i g_1 \frac{\tau_3}{2} \tilde{B}_\mu \Phi_\mathrm{R} - i g_2 \Phi_\mathrm{R} \frac{\tau^a}{2} A^a_{R\mu} \\
D_\mu \Phi_\mathrm{U} &= \partial_\mu \Phi_\mathrm{U} + i g_0 \frac{\tau^a}{2} \tilde{W}^a_\mu \Phi_\mathrm{U} - i g_1 \Phi_\mathrm{U} \frac{\tau_3}{2} \tilde{B}_\mu
\end{aligned} \tag{6}$$

$$\begin{aligned}
D_\mu \Psi_{iL} &= (\partial_\mu + i g_0 \tilde{W}^a_\mu \frac{\tau^a}{2} + \frac{i}{2} g_1 Y' \tilde{B}_\mu) \Psi_{iL} \\
D_\mu \Psi_{iR} &= (\partial_\mu + i g_1 \tilde{B}_\mu \frac{\tau^3}{2} + \frac{i}{2} g_1 Y' \tilde{B}_\mu) \Psi_{iR}
\end{aligned} \tag{7}$$

where g_0, g_1 and g_2 are the coupling constants associated to the groups $SU(2)_L$, $U(1)$ and $SU(2)'_L \otimes SU(2)'_R$ respectively.

For the aim of simplicity, it has been assumed an interchange symmetry between Φ_L and Φ_R in the potential and in the kinetic terms. Writing the scalars in "polar parametrization", *i.e.* $\Phi_\mathrm{L} = \rho_L L$, $\Phi_\mathrm{R} = \rho_R R$ and $\Phi_\mathrm{U} = \rho_U U$ where L, R and U are unitary matrices, the potential gets a simpler form:

$$\begin{aligned}
V(\rho_U, \rho_L, \rho_R) = 2\mu^2 &\left[(\rho_L + u)^2 + (\rho_R + u)^2\right] + \lambda \left[(\rho_L + u)^4 + (\rho_R + u)^4\right] \\
&+ 2m^2 (\rho_U + v)^2 + h(\rho_U + v)^4 + 2 f_3 (\rho_L + u)^2 (\rho_R + u)^2 \\
&+ 2 f (\rho_U + v)^2 \left[(\rho_L + u)^2 + (\rho_R + u)^2\right]
\end{aligned} \tag{8}$$

The scalar fields acquire a vacuum expectation value (vev): $\langle \rho_U \rangle = v$ and $\langle \rho_L \rangle = \langle \rho_R \rangle = u$ which spontaneously break the original gauge symmetry down to $U(1)_\mathrm{em}$.

In the true vacuum, nontrivial mass matrices appear in the scalar and the vector sectors implying that the mass eigenstates are different from the "flavor" ones. In the case of scalars, the relationship between flavor and mass eigenvectors is given, in the limit $u \gg v$, by:

$$\begin{bmatrix} \rho_L \\ \rho_R \\ \rho_U \end{bmatrix} = \begin{bmatrix} \frac{1}{\sqrt{2}} & \frac{1}{\sqrt{2}} - \frac{q^2}{s_\varphi^2} r & -\frac{q}{s_\varphi}\sqrt{r} \\ -\frac{1}{\sqrt{2}} & \frac{1}{\sqrt{2}}(1 - \frac{q^2}{s_\varphi^2} r) & -\frac{q}{s_\varphi}\sqrt{r} \\ 0 & \frac{q}{s_\varphi}\sqrt{2r} & 1 - \frac{q^2}{s_\varphi^2} r \end{bmatrix} \begin{bmatrix} H_L \\ H_R \\ H \end{bmatrix} \tag{9}$$

where the variables r and q are

$$\begin{aligned}
r &= \frac{v^2}{u^2} \frac{g^2}{g_2^2} \qquad q = \frac{f}{f_3 + \lambda}, \\
s_\varphi &= \sin(\varphi) = \frac{g_0}{\sqrt{g_0^2 + g_2^2}}
\end{aligned} \tag{10}$$

and, H_L and H_R are the physical heavy scalar while H denotes the standard-like Higgs boson.

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The mass eigenvalues for the scalar fields, in the limit $u \gg v$, are given by:

$$\begin{aligned} M_H^2 &= 8v^2(h - 2\frac{f^2}{f_3+\lambda}) \\ M_{H_L}^2 &= 8u^2(\lambda - f_3) \\ M_{H_R}^2 &= 8u^2(\lambda + f_3) \end{aligned} \tag{11}$$

Similarly, in the vector sector, the relationship between flavor and mass eigenvectors is given by:

$$\begin{bmatrix} \tilde{W}^{\pm} \\ A_L^{\pm} \end{bmatrix} = \begin{bmatrix} c_\varphi(1 - s_\varphi^2 r) & -s_\varphi(1 + c_\varphi^2 r) \\ s_\varphi(1 + c_\varphi^2 r) & c_\varphi(1 - s_\varphi^2 r) \end{bmatrix} \begin{bmatrix} W^{\pm} \\ V_L^{\pm} \end{bmatrix} \tag{12}$$

$$\begin{bmatrix} \tilde{W}_3 \\ \tilde{B} \\ A_L^3 \\ A_R^3 \end{bmatrix} = \begin{bmatrix} c_\varphi s_\theta & c_\varphi(c_\theta - \frac{s_\varphi^2}{c_\theta} r) & -s_\varphi(1 + c_\varphi^2 r) & \frac{c_\varphi s_\varphi s_\theta^4 \sqrt{P}}{c_\theta^3(1-2c_\theta^2)} r \\ \sqrt{P} & -\frac{s_\theta}{c_\theta}\sqrt{P}(1 - \frac{s_\varphi^2 s_\theta^2}{c_\theta^4} r) & -\frac{c_\varphi s_\varphi s_\theta \sqrt{P}}{1-2c_\theta^2} r & -\frac{s_\varphi s_\theta}{c_\theta}(1 + \frac{s_\theta^2 P}{c_\theta^4} r) \\ s_\varphi s_\theta & s_\varphi c_\theta(1 + \frac{c_\varphi^2}{c_\theta^2} r) & c_\varphi(1 - s_\varphi^2 r) & -\frac{s_\theta^2 P^{3/2}}{c_\theta^3(1-2c_\theta^2)} r \\ s_\varphi s_\theta & -\frac{s_\varphi s_\theta^2}{c_\theta}(1 + \frac{P}{c_\theta^4} r) & \frac{c_\varphi^3 s_\theta^2}{1-2c_\theta^2} r & \frac{\sqrt{P}}{c}{}_\theta(1 - \frac{s_\varphi^2 s_\theta^4}{c_\theta^4} r) \end{bmatrix} \begin{bmatrix} A \\ Z \\ V_L^0 \\ V_R^0 \end{bmatrix} \tag{13}$$

Where, as usual $A, W^{\pm}$ and Z represent the photon and the standard weak gauge bosons while $V_L^{\pm,0}$ and $V_R^{\pm,0}$ are the physical heavy vector bosons and θ represents the Weinberg angle.

The mass eigenvalues of the vector sector are:

$$\begin{aligned} M_A^2 &= 0 \\ M_Z^2 &= \frac{v^2}{4}\frac{g^2}{c_\theta^2}(1 - r s_\varphi^2 \frac{1 - 2c_\theta^2 + 2c_\theta^4}{c_\theta^4} + \cdots) \\ M_{V_L^0}^2 &= \frac{v^2}{4} g^2 (\frac{1}{r c_\varphi^2} + \frac{s_\varphi^2}{c_\varphi^2} - r s_\varphi^2 \frac{c_\theta^2}{1 - 2c_\theta^2} + \cdots) \\ M_{V_R^0}^2 &= \frac{v^2}{4}\frac{g^2}{c_\theta^2}(\frac{1}{r}\frac{c_\theta^4}{P} + \frac{s_\varphi^2 s_\theta^4}{P} + r\frac{s_\varphi^2 s_\theta^8}{c_\theta^4(1 - 2c_\theta^2)} + \cdots) \\ M_{V_R}^2 &= \frac{1}{4} g_2^2 u^2 \\ M_W^2 &= \frac{v^2}{4} g^2 (1 - r s_\varphi^2 + \cdots) \\ M_{V_L}^2 &= \frac{v^2}{4} g^2 (\frac{1}{r}\frac{1}{c_\varphi^2} + \frac{s_\varphi^2}{c_\varphi^2} + r s_\varphi^2 + \cdots) \end{aligned} \tag{14}$$

From the equations above it is easy to see that, when $u \gg v$, it is possible to write the r parameter and c_φ in the convenient form:

$$r \approx \frac{M_W^2}{M_{V_R}^2} \tag{15}$$

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Table 4
Summary of the free parameters of the model.

Parameters	Meaning
u	Scale at which the extended symmetry breaks down
f	Parameter of quartic interactions in the potential
M_{H_L}	Mass of the heavy scalar H_L
M_{H_R}	Mass of the heavy scalar H_R
M_{V_L}	Mass of the heavy vector $V_L^{\pm}, V_L^0$
M_{V_R}	Mass of the heavy vector $V_R^{\pm}, V_R^0$

$$c_\varphi \approx \frac{M_{V_R}}{M_{V_L}} \tag{16}$$

Notice that due to the representation assignments, the fermions only couple to the gauge bosons $\tilde{W}$ and B so their coupling to the heavy vector mass eigenstates arise only through mixing terms.

In this model, as shown in the Lagrangian, Yukawa terms involving only the fermions and the scalar field Φ_{U} are allowed by the global symmetry. The Yukawa Lagrangian can be expanded as follows:

$$\mathcal{L}_Y = \sum_{i.j}^{3}[y_{ij}^d(\bar{L}_i^q \Phi_U)R_j^d + y_{ij}^u(\bar{L}_i^q \tilde{\Phi_U})R_j^u + y_{ij}^{l_d}(\bar{L}_i^l \Phi_U)R_j^{l_d} + y_{ij}^{l_u}(\bar{L}_i^l \Phi_U)R_j^{l_u} + \text{h.c.}] \tag{17}$$

where the components of the scalar field are:

$$\Phi_U = \begin{bmatrix} i w^+ \\ \frac{(v+\rho_{\mathrm{U}})+iz}{\sqrt{2}}, \end{bmatrix} \tag{18}$$

and

$$\rho_U = (1 - \frac{q^2}{s_\varphi^2} r)H + \frac{q}{s_\varphi}\sqrt{2r}H_R. \tag{19}$$

It is important to note that the parameters of the scalar potential have theoretical restrictions, which are described below:

$$\begin{aligned} &\mu^2 < 0 \qquad m^2 < 0 \qquad f > 0 \\ &\lambda - f_3 > 0 \qquad h > f\frac{m^2}{\mu^2} \qquad \lambda + f_3 > 2f\frac{\mu^2}{m^2} \end{aligned} \tag{20}$$

The restrictions for μ^2 and m^2 are imposed so that the potential acquires a vacuum expectation value other than zero. On the other hand the restriction for f comes from the decoupling of the model to the standard model with Higgs. The remaining restrictions in equation (20) derive from the positivity of the mass spectrum of the scalar fields.

Finally, we offer in Table 4 a summary of the free parameters of the model.

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