PROCEEDINGS OF SCIENCE

Non–linear Higgs in a left–right context

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Non–linear operators up to the p^4–order in the Lagrangian expansion are clearly and completely listed for the CP–conserving case in a larger symmetry group $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ and within a non–linear electroweak chiral context coupled to a light dynamical Higgs. The physical effects induced by integrating out the right handed fields from the physical spectrum are briefly commented. Relevant set of effective operators have also been identified at low energies.

Proceedings of the Corfu Summer Institute 2015 "School and Workshops on Elementary Particle Physics and Gravity"
1-27 September 2015
Corfu, Greece

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†Special thanks to Ifigeneia Moraiti and Prof. George Zoupanos for their valuable help during the conference, and for providing us also the best environment along the lectures and seminars.

PoS(CORFU2015)063

1. Effective Lagrangian

The underlying framework relies in a non–linearly realized left–right model coupled to a light Higgs particle. Calling for the larger local group $\mathscr{G} = SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ in an electroweak non–linear σ–model, the Goldstone bosons are parametrized via the dimensionless unitary matrices $\mathbf{U}_L(x)$ and $\mathbf{U}_R(x)$ for the symmetry group $SU(2)_L \times SU(2)_R$, and defined as

$$\mathbf{U}_{L(R)}(x) = e^{i\,\tau_a\,\pi^a_{L(R)}(x)/f_{L(R)}}\,, \tag{1.1}$$

with $\pi^a_{L(R)}(x)$ the corresponding GB fields suppressed by their associated non–linear sigma model scale $f_{L(R)}$. This non–linear effective set–up is coupled a posteriori to a Higgs scalar singlet h through powers of h/f_L [1], via the generic light Higgs polynomial functions $\mathscr{F}(h)$ [2]

$$\mathscr{F}_i(h) \equiv 1 + 2\,a_i\,\frac{h}{f_L} + b_i\,\frac{h^2}{f_L^2} + \mathscr{O}\left(\frac{h^3}{f_L^3}\right). \tag{1.2}$$

Accounting only for the bosonic sector of the model[1], the NP departures with respect to the SM Lagrangian $\mathscr{L}_0$ are encoded through the effective Lagrangian

$$\mathscr{L}_{\rm chiral} = \mathscr{L}_0 + \mathscr{L}_{0,R} + \mathscr{L}_{0,LR} + \Delta\mathscr{L}_{\rm CP} + \Delta\mathscr{L}_{{\rm CP},LR}\,. \tag{1.3}$$

The first two pieces in $\mathscr{L}_{\rm chiral}$ read as

$$\begin{aligned}
\mathscr{L}_0 = &-\frac{1}{4}B_{\mu\nu}B^{\mu\nu} - \frac{1}{4}W^a_{\mu\nu,L}W_L^{\mu\nu,a} - \frac{1}{4}G^a_{\mu\nu}G^{\mu\nu,a} + \\
&+\frac{1}{2}(\partial_\mu h)(\partial^\mu h) - V(h) - \frac{f_L^2}{4}{\rm Tr}\left(\mathbf{V}_L^\mu\mathbf{V}_{\mu,L}\right)\left(1+\frac{h}{f_L}\right)^2 + \\
&+ i\bar{q}_L\not{D}q_L + i\bar{l}_L\not{D}l_L\,,
\end{aligned} \tag{1.4}$$

$$\begin{aligned}
\mathscr{L}_{0,R} = &-\frac{1}{4}W^a_{\mu\nu,R}W_R^{\mu\nu,a} - \frac{f_R^2}{4}{\rm Tr}\left(\mathbf{V}_R^\mu\mathbf{V}_{\mu,R}\right)\left(1+\frac{h}{f_L}\right)^2 + \\
&+ i\bar{q}_R\not{D}q_R + i\bar{l}_R\not{D}l_R\,,
\end{aligned} \tag{1.5}$$

where the adjoints $SU(2)_{L(R)}$–covariant vectorial $\mathbf{V}^\mu_{L(R)}$ and the covariant scalar $\mathbf{T}_{L(R)}$ are defined as

$$\mathbf{V}^\mu_\chi \equiv \left(D^\mu\mathbf{U}_\chi\right)\mathbf{U}_\chi^\dagger\,, \qquad\qquad \mathbf{T}_\chi \equiv \mathbf{U}_\chi\,\tau^3_\chi\,\mathbf{U}^\dagger_\chi\,, \tag{1.6}$$

with $\chi = L,R$, and the corresponding covariant derivative for both of the Goldstone matrices $\mathbf{U}_{L(R)}(x)$ introduced as

$$D^\mu\mathbf{U}_\chi \equiv \partial^\mu\mathbf{U}_\chi + \frac{i}{2}g_\chi W^{\mu,a}_\chi\,\tau^a_\chi\,\mathbf{U}_\chi - \frac{i}{2}g'B^\mu\,\mathbf{U}_\chi\,\tau^3_\chi \tag{1.7}$$

where τ^a_χ stands for the $SU(2)_\chi$ generators ($\chi \equiv L, R$), with the $SU(2)_L$, $SU(2)_R$ and $U(1)_{B-L}$ gauge fields denoted by $W^{a\mu}_L$, $W^{a\mu}_R$ and B^μ correspondingly, and the associated gauge couplings

[1] See [3, 4, 5, 6] for non–linear analysis including fermions.

g_L, g_R and g' respectively. The corresponding $SU(2)_R$–counterparts for the strength gauge kinetic term and the custodial conserving operator at the Lagrangian $\mathscr{L}_0$ are parametrized by $\mathscr{L}_{0,R}$ in (1.5), entailing thus an additional scale f_R that encodes the new high energy scale effects introduced in the scenario once the SM local symmetry group $\mathscr{G}_{SM}$ is extended to $\mathscr{G}$.

The left and right handed structures (LH and RH respectively) in (1.6) can be mixed to build up left-right handed operators (LRH), via the proper insertions of the Goldstone matrices $\mathbf{U}_L$ and $\mathbf{U}_R$, more specifically, through the following definitions [7]

$$\widetilde{\mathbf{V}}^\mu_\chi \equiv \mathbf{U}^\dagger_\chi \mathbf{V}^\mu_\chi \mathbf{U}_\chi, \qquad \widetilde{\mathbf{T}}_\chi \equiv \mathbf{U}^\dagger_\chi \mathbf{T}_\chi \mathbf{U}_\chi, \qquad \widetilde{W}^{\mu\nu}_\chi \equiv \mathbf{U}^\dagger_\chi W^{\mu\nu}_\chi \mathbf{U}_\chi, \tag{1.8}$$

where $W^{\mu\nu}_\chi \equiv W^{\mu\nu,a}_\chi \tau^a_\chi/2$. The p^2-interplaying Lagrangian $\mathscr{L}_{0,LR}$ in (1.3) is thus

$$\mathscr{L}_{0,LR} = -\frac{1}{2}\,\mathrm{Tr}\left(\widetilde{W}^{\mu\nu}_L \widetilde{W}_{\mu\nu,R}\right) - \frac{f_L f_R}{2}\,\mathrm{Tr}\left(\widetilde{\mathbf{V}}^\mu_L \widetilde{\mathbf{V}}_{\mu,R}\right)\left(1+\frac{h}{f_L}\right)^2. \tag{1.9}$$

Non-zero mass mixing terms among the LH and RH gauge fields are triggered by the latter Lagrangian in the unitary gauge, leading then to diagonalize the gauge sector in order to obtain the required physical gauge masses [8].

1.1 CP-preserving deviations: $\Delta\mathscr{L}_{\mathrm{CP}}$

Non–zero NP departures with respect to those described in $\mathscr{L}_0 + \mathscr{L}_{0,R} + \mathscr{L}_{0,LR}$ are parametrized through the remaining last two pieces in (1.3): $\Delta\mathscr{L}_{\mathrm{CP}}$ and $\Delta\mathscr{L}_{\mathrm{CP},LR}$. The former contains LH and RH covariant objects up to the p^4–order as

$$\Delta\mathscr{L}_{\mathrm{CP}} = \Delta\mathscr{L}_{\mathrm{CP},L} + \Delta\mathscr{L}_{\mathrm{CP},R} \tag{1.10}$$

where each one of the components are written down as

$$\Delta\mathscr{L}_{\mathrm{CP},L} = c_B\,\mathscr{P}_B(h) \;+\; \sum_i c_{i,L}\,\mathscr{P}_{i,L}(h) \;+\; \sum_{j=1}^{26} c_{j,L}\,\mathscr{P}_{j,L}(h)\,, \tag{1.11}$$

$$\Delta\mathscr{L}_{\mathrm{CP},R} = \sum_i c_{i,R}\,\mathscr{P}_{i,R}(h) \;+\; \sum_{j=1}^{26} c_{j,R}\,\mathscr{P}_{j,R}(h)\,. \tag{1.12}$$

where the summation over the index i runs over the labels $i = \{W,C,T\}$. The model–dependent constant coefficients c_B, $c_{i,L}$ and $c_{i,R}$ are denoting correspondingly the weighting coefficients for the LH and RH operators, whilst the first two terms of $\Delta\mathscr{L}_{\mathrm{CP},L}$ in (1.11) and the first term in (1.12) can be jointly written as

$$\begin{aligned}
\mathscr{P}_B(h) &= -\frac{g'^2}{4}\,B_{\mu\nu}B^{\mu\nu}\,\mathscr{F}_B(h)\,,\\
\mathscr{P}_{W,\chi}(h) &= -\frac{g_\chi^2}{4}\,W^a_{\mu\nu,\chi}W^{\mu\nu,a}_\chi\,\mathscr{F}_{W,\chi}(h)\,,\\
\mathscr{P}_{C,\chi}(h) &= -\frac{f_\chi^2}{4}\mathrm{Tr}\left(\mathbf{V}^\mu_\chi \mathbf{V}_{\mu,\chi}\right)\mathscr{F}_{C,\chi}(h)\,,\\
\mathscr{P}_{T,\chi}(h) &= \frac{f_\chi^2}{4}\left(\mathrm{Tr}\left(\mathbf{T}_\chi \mathbf{V}^\mu_\chi\right)\right)^2\mathscr{F}_{T,\chi}(h)\,,
\end{aligned} \tag{1.13}$$

with suffix χ labelling again as $\chi = L, R$. The generic function $\mathcal{F}_i(h)$ follows definition (1.2). No gluonic operator has been included.

The complete linearly independent set of 26 CP–conserving non–linear operators $\mathcal{P}_{i,L}(h)$ (first term in the second line of $\Delta\mathcal{L}_{\mathrm{CP},L}$, Eq. (1.11)) have completely been listed in Refs. [2, 9], whereas the symmetric right handed counterpart, described by the set of 26 operators $\mathcal{P}_{i,R}(h)$ (second term in $\Delta\mathcal{L}_{\mathrm{CP},R}$ of Eq. (1.12)) was listed in [7, 8]. A bunch of the whole set $\{\mathcal{P}_{i,L}(h),\,\mathcal{P}_{i,R}(h)\}$ is given by

$$
\begin{aligned}
\mathcal{P}_{1,\chi}(h) &= g_\chi\, g'\, B_{\mu\nu}\,\mathrm{Tr}\Big(\mathbf{T}_\chi\, W_\chi^{\mu\nu}\Big)\,\mathcal{F}_{1,\chi}(h)\,,\\
\mathcal{P}_{2,\chi}(h) &= i\,g'\, B_{\mu\nu}\,\mathrm{Tr}\Big(\mathbf{T}_\chi\,\Big[\mathbf{V}_\chi^\mu,\mathbf{V}_\chi^\nu\Big]\Big)\,\mathcal{F}_{2,\chi}(h)\,,\\
\mathcal{P}_{3,\chi}(h) &= i\,g_\chi\,\mathrm{Tr}\Big(W_\chi^{\mu\nu}\,\Big[\mathbf{V}_{\mu,\chi},\mathbf{V}_{\nu,\chi}\Big]\Big)\,\mathcal{F}_{3,\chi}(h)\,,\\
&\;\;\vdots\\
\mathcal{P}_{25,\chi}(h) &= \Big(\mathrm{Tr}\Big(\mathbf{T}_\chi\,\mathbf{V}_\chi^\mu\Big)\Big)^2\,\partial_\nu\partial^\nu\mathcal{F}_{25,\chi}(h)\,,\\
\mathcal{P}_{26,\chi}(h) &= \Big(\mathrm{Tr}\Big(\mathbf{T}_\chi\,\mathbf{V}_\chi^\mu\Big)\,\mathrm{Tr}\Big(\mathbf{T}_\chi\,\mathbf{V}_\chi^\nu\Big)\Big)^2\,\mathcal{F}_{26,\chi}(h)\,.
\end{aligned}
\tag{1.14}
$$

The covariant derivative $\mathcal{D}_\mu$ in Eq. (1.14) acts on a field transforming in the adjoint representation of $SU(2)_L$, being defined as

$$
\mathcal{D}^\mu\mathbf{V}_\chi^\nu \equiv \partial^\mu\mathbf{V}_\chi^\nu + i\,g_\chi\Big[W_\chi^\mu,\mathbf{V}_\chi^\nu\Big]\,,\qquad \chi = L, R\,. \tag{1.15}
$$

The contribution $\Delta\mathcal{L}_{\mathrm{CP},L}$ has already been provided in [2, 9] in the context of purely EW chiral effective theories coupled to a light Higgs, whereas part of $\Delta\mathcal{L}_{\mathrm{CP},L}$ and $\Delta\mathcal{L}_{\mathrm{CP},R}$ were partially analysed for the left–right symmetric frameworks in [10, 11], and finally completed in [7]. See [7, 8] for more details.

1.2 Left–right interplaying departures: $\Delta\mathcal{L}_{\mathrm{CP},LR}$

Finally, $\Delta\mathcal{L}_{\mathrm{CP},LR}$ parametrizes those mixing interacting terms among the $SU(2)_L$ and $SU(2)_R$–covariant objects up to the p^4–order in the Lagrangian expansion, permitted by the underlying left–right symmetry, and encoded through

$$
\Delta\mathcal{L}_{\mathrm{CP},LR} \;=\; \sum_k c_{k,LR}\,\mathcal{P}_{k,LR}(h) \;+\; \sum_{i=2,\,i\neq 4}^{26} c_{i(j),LR}\,\mathcal{P}_{i(j),LR}(h)\,, \tag{1.16}
$$

where the summation over the index k runs again over the labels $k = \{W, C, T\}$. The index j spans over all the possible operators that can be built up from the set of 26 operators $\mathcal{P}_{i,\chi}(h)$ in (1.11)–(1.12), and here labelled as $\mathcal{P}_{i(j),LR}(h)$ together with their corresponding coefficients $c_{i(j)}$. The first term in $\Delta\mathscr{L}_{\mathrm{CP},LR}$ encodes the non-linear mixing operators

$$
\begin{aligned}
\mathscr{P}_{W,LR}(h) &= -\frac{1}{2}\,g_L\,g_R\,\mathrm{Tr}\left(\widetilde{W}_L^{\mu\nu}\,\widetilde{W}_{\mu\nu,R}\right)\,\mathscr{F}_{W,LR}(h)\,,\\
\mathscr{P}_{C,LR}(h) &= \frac{1}{2}\,f_L\,f_R\,\mathrm{Tr}\left(\widetilde{\mathbf{V}}_L^{\mu}\widetilde{\mathbf{V}}_{\mu,R}\right)\,\mathscr{F}_{C,LR}(h)\,,\\
\mathscr{P}_{T,LR}(h) &= \frac{1}{2}\,f_L\,f_R\,\mathrm{Tr}\left(\widetilde{\mathbf{T}}_L\,\widetilde{\mathbf{V}}_L^{\mu}\right)\,\mathrm{Tr}\left(\widetilde{\mathbf{T}}_R\,\widetilde{\mathbf{V}}_{\mu,R}\right)\mathscr{F}_{T,LR}(h)\,.
\end{aligned}
\tag{1.17}
$$

A bunch of the whole set of 75 linearly independent operators $\mathscr{P}_{i(j),LR}(h)$ (2nd second term of $\Delta\mathscr{L}_{\mathrm{CP},LR}$, reported in [8, 12]) is provided again for briefness reasons:

$$
\begin{aligned}
&\mathscr{P}_{2(1)}(h) = i\,g'\,B_{\mu\nu}\,\mathrm{Tr}\left(\widetilde{\mathbf{T}}_L\left[\widetilde{\mathbf{V}}_L^{\mu},\widetilde{\mathbf{V}}_R^{\nu}\right]\right)\,\mathscr{F}_{2(1)}(h)\,,\\
&\mathscr{P}_{3(1)}(h) = i\,g_L\,\mathrm{Tr}\left(\widetilde{W}_L^{\mu\nu}\left[\widetilde{\mathbf{V}}_{\mu,R},\widetilde{\mathbf{V}}_{\nu,R}\right]\right)\,\mathscr{F}_{3(1)}(h)\,,\\
&\mathscr{P}_{3(2)}(h) = i\,g_R\,\mathrm{Tr}\left(\widetilde{W}_R^{\mu\nu}\left[\widetilde{\mathbf{V}}_{\mu,L},\widetilde{\mathbf{V}}_{\nu,L}\right]\right)\,\mathscr{F}_{3(2)}(h)\,,\\
&\mathscr{P}_{3(3)}(h) = i\,g_L\,\mathrm{Tr}\left(\widetilde{W}_L^{\mu\nu}\left[\widetilde{\mathbf{V}}_{\mu,L},\widetilde{\mathbf{V}}_{\nu,R}\right]\right)\,\mathscr{F}_{3(3)}(h)\,,\\
&\qquad\vdots\\
&\mathscr{P}_{7(1)}(h) = \mathrm{Tr}\left(\widetilde{\mathbf{V}}_L^{\mu}\,\widetilde{\mathbf{V}}_{\mu,R}\right)\,\partial_\nu\partial^\nu\mathscr{F}_{7(1)}(h)\,,\\
&\mathscr{P}_{8(1)}(h) = \mathrm{Tr}\left(\widetilde{\mathbf{V}}_L^{\mu}\,\widetilde{\mathbf{V}}_R^{\nu}\right)\,\partial_\mu\mathscr{F}_{8(1)}(h)\,\partial_\nu\mathscr{F}'_{8(1)}(h)\,,\\
&\mathscr{P}_{9(1)}(h) = \mathrm{Tr}\left(\mathscr{D}_\mu\widetilde{\mathbf{V}}_L^{\mu}\,\mathscr{D}_\nu\widetilde{\mathbf{V}}_R^{\nu}\right)\,\mathscr{F}_{9(1)}(h)\,,\\
&\mathscr{P}_{10(1)}(h) = \mathrm{Tr}\left(\widetilde{\mathbf{V}}_L^{\nu}\,\mathscr{D}_\mu\widetilde{\mathbf{V}}_R^{\mu}\right)\,\partial_\nu\mathscr{F}_{10(1)}(h)\,,\\
&\mathscr{P}_{10(2)}(h) = \mathrm{Tr}\left(\widetilde{\mathbf{V}}_R^{\nu}\,\mathscr{D}_\mu\widetilde{\mathbf{V}}_L^{\mu}\right)\,\partial_\nu\mathscr{F}_{10(2)}(h)\,,\\
&\qquad\vdots\\
&\mathscr{P}_{26(1)}(h) = \left(\mathrm{Tr}\left(\widetilde{\mathbf{T}}_L\,\widetilde{\mathbf{V}}_L^{\mu}\right)\mathrm{Tr}\left(\widetilde{\mathbf{T}}_R\,\widetilde{\mathbf{V}}_{\mu,R}\right)\right)^2\mathscr{F}_{26(1)}(h)\,,\\
&\mathscr{P}_{26(2)}(h) = \left(\mathrm{Tr}\left(\widetilde{\mathbf{T}}_L\,\widetilde{\mathbf{V}}_L^{\mu}\right)\mathrm{Tr}\left(\widetilde{\mathbf{T}}_R\,\widetilde{\mathbf{V}}_R^{\nu}\right)\right)^2\mathscr{F}_{26(2)}(h)\,,\\
&\mathscr{P}_{26(3)}(h) = \mathrm{Tr}\left(\widetilde{\mathbf{T}}_L\,\widetilde{\mathbf{V}}_L^{\mu}\right)\mathrm{Tr}\left(\widetilde{\mathbf{T}}_R\,\widetilde{\mathbf{V}}_{\mu,R}\right)\left(\mathrm{Tr}\left(\widetilde{\mathbf{T}}_L\,\widetilde{\mathbf{V}}_L^{\nu}\right)\right)^2\mathscr{F}_{26(3)}(h)\,,\\
&\mathscr{P}_{26(4)}(h) = \mathrm{Tr}\left(\widetilde{\mathbf{T}}_L\,\widetilde{\mathbf{V}}_L^{\mu}\right)\mathrm{Tr}\left(\widetilde{\mathbf{T}}_R\,\widetilde{\mathbf{V}}_{\mu,R}\right)\left(\mathrm{Tr}\left(\widetilde{\mathbf{T}}_R\,\widetilde{\mathbf{V}}_R^{\nu}\right)\right)^2\mathscr{F}_{26(4)}(h)\,,
\end{aligned}
\tag{1.18}
$$

where the suffix LR in all $\mathscr{P}_{i(j),LR}(h)$ and their corresponding $\mathscr{F}_{i(j),LR}(h)$ has been omitted as well in (1.18). The complete set of 75 linearly independent operators $\mathscr{P}_{i(j),LR}(h)$ were listed in [8, 12]. Among them, 23 operators were missing in the left-right symmetric EW chiral treatment of

Refs. [10, 11] (see [8, 12] for more details). The corresponding CP–violating counterparts of $\Delta\mathscr{L}_{\rm CP}$ and $\Delta\mathscr{L}_{{\rm CP},LR}$ have been completely listed and studied in [16].

At the unitary gauge, non-zero mass mixing terms among the LH and RH gauge fields are triggered by the operator $\mathscr{P}_{C,LR}(h)$, for both the charged gauge basis $\{W^\pm_{\mu,L}, W^\pm_{\mu,R}\}$ and the neutral one $\{W^3_{\mu,L}, W^3_{\mu,R}, B_\mu\}$. A rotation in the gauge sector is in order to obtain the required gauge boson masses in the physical basis $\{W^\pm_\mu, W'^\pm_\mu\}$ and $\{A_\mu, Z_\mu, Z'_\mu\}$ respectively. A mass range for the gauge field Z' is predicted in terms of the W'-mass and the gauge coupling g_R [8, 12]. Interpreting the observed excesses at the ATLAS Collaborations in the WZ–final state, and by the CMS Collaboration in the e^+e^-jj, Wh and jj–final states, to be induced by decays of a heavy boson W' in the 1.8–2 TeV mass range, and assuming the coupling g_R in the range $g_R \approx 0.45-0.6$ as determined in [13], it is possible to predict the mass range $2.4\,\text{TeV} < M_{Z'} < 4\,\text{TeV}$ [8, 12]. In addition, a $M_{W'}$–range of 1.8–2 TeV entails a scale $f_R \sim$ 6–8 TeV [8]. A more detailed interpretation of the diboson excess via a left–right non-linear Higgs approach is done in [12].

Generally, a higher energy scale f_R points towards higher masses $M_{W'}$ and $M_{Z'}$, implying a vanishing mixing angle among the charged gauge fields $W^\pm_{\mu,L}$ and $W^\pm_{\mu,R}$, neither a mixing among the set of neutral fields $\{W^3_{\mu,L}, B_\mu\}$ with the field $W^3_{\mu,R}$. RH gauge fields can thus be directly linked to the eigenstate basis as $W^\pm_{\mu,R} = W'^\pm_\mu$ and $W^3_{\mu,R} = -Z'_\mu$ for a higher energy scale f_R. Heavy right handed gauge fields can thus be integrated out from the physical spectrum of the model, triggering therefore physical effects that will be manifested at lower energies in the effective Lagrangian.

1.3 Integrating-out heavy right handed fields

Through the equations of motion for the gauge and Higgs fields, RH gauge fields can be integrated out from the physical spectrum via the relations [7, 8]

$$\mathbf{V}^\mu_R \equiv -\varepsilon\,\mathbf{V}^\mu_L\,, \qquad \text{with} \qquad \varepsilon \equiv \frac{f_L}{f_R}\,(1+c_{C,LR}) \tag{1.19}$$

that can be translated into the unitary gauge as

$$W^\pm_{\mu,R} \quad\Rightarrow\quad -\frac{g_L}{g_R}\,\varepsilon\,W^\pm_{\mu,L}\,, \qquad\qquad W^3_{\mu,R} \quad\Rightarrow\quad \frac{g'}{g_R}\,(1+\varepsilon)\,B_\mu - \frac{g_L}{g_R}\,\varepsilon\,W^3_{\mu,L}\,. \tag{1.20}$$

All the RH and LRH operators will collapse onto the LH ones after plugging back the Eq. (1.19) through (1.12)-(1.14) (for $\chi = R$) and (1.16)-(1.18), affecting thus the corresponding global coefficients $c_{i,L}$ in a generic manner as

$$c_{i,L} \;\Longrightarrow\; \tilde{c}_{i,L} = c_{i,L} + \sum_{k=1}^{4} \varepsilon^k\,\mathscr{F}^{(k)}\left(c_{i,R},\, c_{i(j)},\, c_{l(m)}\right) \tag{1.21}$$

where the functions $\mathscr{F}^{(k)}\left(c_{i,R}, c_{i(j)}, c_{l(m)}\right)$ will encode linear combinations on the coefficients $c_{i,R}$, $c_{i(j)}$ and additional mixing left-right operators via $c_{l(m)}$ (see Ref. [8] for more details). The contributions induced onto the left handed operators are suppressed by powers of the ratio f_L/f_R, being determined by the number of fields $\mathbf{V}^\mu_R$ through each one of the right and left–right operators. Consequently, in the limiting case $f_L \ll f_R$ at low energies, it is realized that the set of non-linear operators $\{\mathscr{P}_B, \mathscr{P}_{C,L}, \mathscr{P}_{T,L}, \mathscr{P}_{1,L}, \mathscr{P}_{2,L}, \mathscr{P}_{4,L}\}$ is sensitive to the contributions, up to the order $\mathscr{O}(\varepsilon)$, from the right handed operators

$$\{\mathscr{P}_{C,R}, \mathscr{P}_{T,R}, \mathscr{P}_{W,R}, \mathscr{P}_{1,R}, \mathscr{P}_{12,R}\} \tag{1.22}$$

and the mixing left–right set

$$\{\mathscr{P}_{C,LR}, \mathscr{P}_{T,LR}, \mathscr{P}_{W,LR}, \mathscr{P}_{3(2)}, \mathscr{P}_{12(1)}, \mathscr{P}_{13(2)}, \mathscr{P}_{17(2)}\}\,. \tag{1.23}$$

This is relevant for the EWPT parameters S and T, as they are sensitive to the effects from $\mathscr{P}_{1,L}$ and $\mathscr{P}_{T,L}$ respectively [8]. The tree-level contributions to the oblique parameters S and T [14] turn out to be

$$\alpha_{\rm em}\,\Delta S = 2\,s_{2W}\,\alpha_{WB} - 8\,e^2\,\tilde{c}_{1,L}\,, \qquad\qquad \alpha_{\rm em}\,\Delta T = 2\,\tilde{c}_{T,L}\,, \tag{1.24}$$

with $\alpha_{\rm em}$ the fine structure constant and the notation $s_{2W} \equiv \sin(2\,\theta_W)$. The coefficient α_{WB} and the redefined ones $\tilde{c}_{1,L}$ and $\tilde{c}_{T,L}$ are defined as

$$\alpha_{WB} \equiv \frac{g'}{2g_R}\left(1-2\frac{g_L}{g_R}\,\varepsilon\right)(1+\varepsilon)\,, \qquad \tilde{c}_{1,L} = c_{1,L} - \frac{1}{4}c_{W,LR} + c_{12(1)}\,, \qquad \tilde{c}_{T,L} = c_{T,L} + c_{T,R} - 2c_{T,LR}\,. \tag{1.25}$$

Furthermore, the triple gauge–boson couplings (TGC) are also sensitive to the induced effects by integrating out the right handed fields, being generically described via [15]

$$\begin{aligned}\frac{\mathscr{L}_{\rm TGV}}{g_{WWV}} = i\Big\{ & g_1^V\left(W^+_{\mu\nu}W^{-\mu}V^\nu - W^+_\mu V_\nu W^{-\mu\nu}\right) + \kappa_V W^+_\mu W^-_\nu V^{\mu\nu} + \\ & - i g_5^V \lambda^{\mu\nu\rho\sigma}\left(W^+_\mu \partial_\rho W^-_\nu - W^-_\nu \partial_\rho W^+_\mu\right) V_\sigma + g_6^V\left(\partial_\mu W^{+\mu}W^{-\nu} - \partial_\mu W^{-\mu}W^{+\nu}\right)V_\nu \Big\}\,,\end{aligned} \tag{1.26}$$

where $V \equiv \{\gamma, Z\}$ and $g_{WW\gamma} \equiv e$, $g_{WWZ} \equiv e\,c_W/s_W$, with $W^\pm_{\mu\nu}$ and $V_{\mu\nu}$ standing for the kinetic part of the implied gauge field strengths. Compact notation $c_W \equiv \cos\theta_W$ and $s_W \equiv \sin\theta_W$ is implicit. Electromagnetic gauge invariance requires $g_1^\gamma = 1$ and $g_5^\gamma = 0$, in consequence the CP-even TGC encoded in (1.26) depends in all generality on six dimensionless couplings g_1^Z, g_5^Z, $g_6^{\gamma,Z}$ and $\kappa_{\gamma,Z}$. Their SM values are $g_1^Z = \kappa_\gamma = \kappa_Z = 1$ and $g_5^Z = g_6^\gamma = g_6^Z = 0$. Additionally, the couplings g_6^V have been introduced to account for the contributions associated to the operators containing the contraction $\mathscr{D}_\mu \mathbf{V}^\mu_L$, with its corresponding $\partial_\mu \mathbf{V}^\mu_L$–part vanishing only for on-shell gauge bosons. When fermion masses are neglected, such contraction can be disregarded. The set of TGC parametrized through $\mathscr{L}_{\rm TGV}$ in (1.26) are written up to the $\mathscr{O}(\varepsilon)$-contributions as

$$g_1^Z = 1 - \frac{2s_W^4}{c_{2W}s_{2W}}\,\alpha_{WB} + \frac{1}{2c_{2W}}\left[\tilde{c}_{T,L} - 4e^2\left(c_{12,L} - \frac{s_W^2\tilde{c}_{1,L}}{c_W^2}\right)\right] - \frac{4e^2 c_{3,L}}{s_{2W}^2}\,,$$

$$\kappa_\gamma = 1 + \frac{c_W}{s_W}\,\alpha_{WB} - \frac{e^2}{s_W^2}\left(2\tilde{c}_{1,L} + 2\tilde{c}_{2,L} + c_{3,L} + 4c_{12,L} + 2c_{13,L}\right)\,,$$

$$\kappa_Z = 1 - \frac{s_{2W}}{2c_{2W}}\,\alpha_{WB} + \frac{\tilde{c}_{T,L}}{2\,c_{2W}} + e^2\left(\frac{2\tilde{c}_{1,L}}{c_{2W}} + \frac{2\tilde{c}_{2,L}}{c_W^2}\right) - \frac{e^2}{s_W^2}\left[\left(\frac{1}{c_{2W}} + 3\right)c_{12,L} + c_{3,L} + 2c_{13,L}\right]\,,$$

$$g_5^Z = -\frac{4e^2}{s_{2W}^2}\,c_{14,L}\,, \qquad g_6^\gamma = \frac{e^2}{s_W^2}\,c_{9,L}\,, \qquad g_6^Z = e^2\left(\frac{4c_{16,L}}{s_{2W}^2} - \frac{c_{9,L}}{c_W^2}\right)\,, \tag{1.27}$$

with

$$\tilde{c}_{2,L} = c_{2,L} + \frac{1}{2}\left(2c_{13(2)} + c_{3(2)}\right)\,. \tag{1.28}$$

Likewise, some pair gauge bosons–Higgs couplings will be affected too. In fact, the vertexes $\{F_{\mu\nu}F^{\mu\nu}h,\, Z_{\mu\nu}Z^{\mu\nu}h,\, F_{\mu\nu}Z^{\mu\nu}h,\, Z_\mu Z^{\mu\nu}\,\partial_\nu h,\, Z_\mu F^{\mu\nu}\,\partial_\nu h\}$, and $\{W_\mu^\dagger W^\mu h,\, Z_\mu Z^\mu h\}$ will depend of linear combinations of the operators $\{\mathscr{P}_B,\, \mathscr{P}_{C,L},\, \mathscr{P}_{T,L},\, \mathscr{P}_{1,L},\, \mathscr{P}_{2,L},\, \mathscr{P}_{4,L}\}$. See [8] for further details on the implied phenomenology and the allowed ranges for the involved operator coefficients.

Finally, by disregarding: i) LH operators with negligible physical impact and irrelevant for the non-linear realization of the dynamics [9], ii) operators redundant for the massless fermion case via EOM, iii) and those ones not directly contributing to any of the couplings listed previously, one can finally disregard 19 LH operators in total [8], remaining thus with an effective set of 17 LH ops. = 31 (set in (1.13) + (1.14)) - 14. From all these considerations it is concluded that a RH gauge sector far above the EW scale will imply a hierarchical case with NP effects parametrized via a much smaller operator basis as the f_L/f_R–suppression would entail, and leaving us therefore with 29 operators in total = 17 LH ops. + 5 RH ops. (in (1.22)) + 7 LRH ops (in (1.23)).

2. Conclusions

Effective Lagrangian approaches are in order to parametrize possible new physical effects detectable at low energies. Concerning only the bosonic gauge sector, we assume here a NP field content pictured by spin–1 resonances from larger local group $\mathscr{G} = SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$, here described via a non–linear EW scenario with a light dynamical Higgs, and up to the p^4-contributions in the Lagrangian expansion.

The analysis provided in this work may also be considered as a generic UV completion of the low energy non–linear treatments of Refs. [23, 24, 25, 26, 27] and Refs. [2, 9, 17], as long as the extended gauge field sector arises out from an energy regime higher than the EW scale. The physical effects induced by integrating out the right handed fields from the physical spectrum are analysed. The relevant set of operators have been identified at low energies, 24 operators in total = 12 left ops. + 5 right ops. (in (1.22)) + 7 left–right ops (in (1.23)). More low energy effects from a higher energy gauge sector [8, 12, 16] could unveil the underlying NP playing a role in our nature, and likely will point towards a better understanding on the origin of the electroweak symmetry breaking mechanism.

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