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Anomalous Weak Boson Couplings: Suggestions from Unitarity and Dynamics †

G.J. Gounaris[a,b], F.M. Renard[c] and G. Tsirigoti[b]

[a]Theory Division, CERN,
CH-1211 Geneva 23, Switzerland,

[b]Department of Theoretical Physics, University of Thessaloniki,
Gr-54006, Thessaloniki, Greece,

[c]Physique Mathématique et Théorique, CNRS-URA 768,
Université de Montpellier II, F-34095 Montpellier Cedex 5.

Abstract

Taking into account the constraints from LEP1 and lower energy experiments, we identify the seven $SU(2) \times U(1)$ gauge invariant purely bosonic $dim = 6$ operators which provide a quite general description of how New Physics could reflect in the bosonic world, if it happens that all new degrees of freedom are too heavy to be directly produced in the future colliders. Five of these operators are CP conserving and the remaining ones are CP violating. We derive the unitarity constraints for the CP violating operators and compare them with the already known constraints for the CP conserving ones. Dynamical renormalizable models are also presented, which partly elucidate what the appearance of each of these operators can teach us on the mechanism of spontaneous gauge symmetry breaking.

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1 Introduction

The intense experimental effort to find traces of New Physics (NP) beyond the Standard Model (SM) has so far given only very weak hints [1, 2]. No new particles, possibly associated with NP have ever been seen in our present accelerators. Moreover, the interactions among the gauge bosons and the light fermions have been thoroughly scrutinized at LEP1 and lower energies, and were found to be fully consistent with the SM predictions. The only slight experimental hints for something beyond the SM that exist at present consist in the well known peculiarities observed in $Z \to b\bar{b}$ [1, 2], and the persisting indications favouring the possibility of some non-vanishing neutrino masses and a huge amount of dark matter in the Universe.

Thus, before the excitation of new particles will become possible, hopefully in one of the contemplated future accelerators, it seems that our main hope to detect hints of NP is by carefully searching for anomalies in interactions among the gauge bosons, the Higgs and the quarks of the third family, since these interactions have not yet been tested to the same level of accuracy as the light fermionic ones [1, 2, 3]. Of course, if the Higgs particle turns out to be above the TeV scale, it will itself be part of NP, inducing new strong interactions mainly among the longitudinal gauge bosons [4]. Although this is a viable possibility, we assume below that it will not be the case in Nature, and that the Higgs will be discovered some day in the mass range of the electroweak breaking scale $v = 1/(\sqrt{2}G_\mu)^{1/2}$.

We therefore contemplate an NP scenario according to which the usual SM Higgs particle exists and is, in a sense, part of the "old" physics. Moreover, in this scenario the NP scale Λ_{NP}, which determines the masses of all new particles, is assumed to be very large. Under such conditions, a quite general way of parametrizing NP is achieved by establishing an effective Lagrangian containing contributions from all possible $SU(3) \times SU(2) \times U(1)$ gauge invariant operators constructed from scalar and gauge bosons fields, as well as the quarks of the third family (together with the gluons). Since the contributions of these operators are scaled by inverse powers of Λ_{NP}, it is plausible to expect that for a sufficiently large NP scale, the $dim = 6$ operators should give in general an adequate description [5]. Of course, it is quite possible that there exist NP aspects whose scale is not really very large, such as the case of a moderately heavy vector boson which mixes with W or Z [6, 7]. In such a case our approximation to retain only $dim = 6$ operators might not be sufficient, and we would have to include in our expansion also higher dimensional operators. In the following we will assume that this is not the case, though.

The complete list of purely bosonic such $dim = 6$ NP operators has been known for some time[1] [8, 9], and recently we have also established the CP invariant operators containing the $(t\ ,b)$ quarks [3]. In the present paper we focus on the purely bosonic operators though, eleven of which conserve the CP symmetry, while the remaining five ones violate it. We next take into account the fact that seven of these purely bosonic operators, (four CP conserving and three violating ones) are already excluded by LEP1 and low energy measurements [8, 10, 11], while another two are completely insensitive to

[1] These operators involve only weak gauge bosons and Higgs.

any conceivable experiments [8]. Discarding all these irrelevant interactions, we conclude that it should be sufficient to describe the purely bosonic part of NP below Λ_{NP}, in terms of an effective Lagrangian which is a linear combination of the seven operators called $\mathcal{O}_W$, $\mathcal{O}_{W\Phi}$, $\mathcal{O}_{B\Phi}$, $\mathcal{O}_{UW}$, $\mathcal{O}_{UB}$, $\overline{\mathcal{O}}_{UW}$, $\overline{\mathcal{O}}_{UB}$. Presently existing LEP1 experiments put only very moderate constraints on these operators [8, 10]. The general aim of the present paper is to give an orientation on the magnitude of the couplings of these operators, based on considerations on the unitarity constraints and on a class of dynamical scenarios.

The first three of these operators, namely $\mathcal{O}_W$, $\mathcal{O}_{W\Phi}$ and $\mathcal{O}_{B\Phi}$, are the only ones involving triple gauge boson couplings. These operators are also CP symmetric, and are the only ones to give anomalous contributions to the process $e^+e^- \to W^+W^-$, which will be studied at LEP2 and NLC [12, 13].

The remaining four operators only induce anomalous $H\gamma\gamma$, $H\gamma Z$, HZZ and HWW couplings, and contain no triple gauge boson vertices [9, 14]. The operators $\mathcal{O}_{UW}$ and $\mathcal{O}_{UB}$ are CP symmetric, while $\overline{\mathcal{O}}_{UW}$, $\overline{\mathcal{O}}_{UB}$ are CP violating. If H is within the LEP2 range, these couplings could be studied there by carefully analysing Higgs-strahlung [9, 14, 15]. Thus, immediately after the "hoped for" discovery of the Higgs, the need to search for its anomalous couplings will arise.

In order to study the NP signatures described by the above operators, it is very useful to first establish the unitarity constraints on their couplings. Such unitarity constraints give relations between the strength of these couplings and the energy scale where either unitarity will be saturated, or (as happened in the old Fermi theory) some of the new degrees of freedom of NP will start being excited. At the technical level such relations are very helpful, since they roughly determine the energies and couplings for which the perturbative results are reliable. In previous works we have established the unitarity constraints for the five CP conserving operators [16]. The first aim of the present paper is to complete this study by giving the unitarity constraints also for the CP violating ones $\overline{\mathcal{O}}_{UW}$, $\overline{\mathcal{O}}_{UB}$. We then summarize the implications from the unitarity relations for all seven operators. These are the "suggestions from Unitarity" alluded to in the title.

The second aim of the present work is to offer examples of dynamical models containing new heavy degrees of freedom, which, after they are integrated out, lead to an NP description in terms of the purely bosonic $dim = 6$ operators. These examples are generalizations of previous ones given for the case of NP operators respecting custodial $SU(2)_c$ symmetry [17]. The usefulness of such examples consists in the fact that they provide a feeling on what type of anomalous interactions could be induced by various kinds of new degrees of freedom. From these examples, we infer that Higgs dependent operators $\mathcal{O}_{UW}$, $\mathcal{O}_{UB}$, $\overline{\mathcal{O}}_{UW}$, $\overline{\mathcal{O}}_{UB}$, if they happen to be created in the model, have their couplings determined by the arbitrary Yukawa type interactions in new physics. Thus, their couplings are not constrained by existing experiments and have a chance to be observable at LEP2 [9], [15]. On the contrary, the purely gauge dependent operator $\mathcal{O}_W$ has its strength determined exclusively by the group properties of the NP particles we have integrated out. We would expect, therefore, that $\mathcal{O}_W$ could only become appreciable if a non-perturbative mechanism enhances it. The same is also true for the operators $\mathcal{O}_{UW}$ and $\mathcal{O}_{UB}$, which were never generated in these models.

Disregarding all contributions which are either unobservable or are very strongly constrained by existing experiments, and restricting to $dim = 6$ operators [8, 10, 11], the effective Lagrangian describing the purely bosonic part of NP at the scale Λ_{NP} is given by

$$\begin{aligned}\mathcal{L}_{NP} &= \lambda_W \frac{g}{M_W^2} \mathcal{O}_W + \frac{f_B g'}{2M_W^2} \mathcal{O}_{B\Phi} + \frac{f_W g}{2M_W^2} \mathcal{O}_{W\Phi} \quad + \\ & d\, \mathcal{O}_{UW} + \frac{d_B}{4}\, \mathcal{O}_{UB} + \overline{d}\, \overline{\mathcal{O}}_{UW} + \frac{\overline{d}_B}{4}\, \overline{\mathcal{O}}_{UB} \qquad , \end{aligned} \tag{1}$$

where

$$\mathcal{O}_{W\Phi} = i\,(D_\mu \Phi)^\dagger \overrightarrow{\tau} \cdot \overrightarrow{W}^{\mu\nu} (D_\nu \Phi) \quad , \tag{2}$$

$$\mathcal{O}_{B\Phi} = i\,(D_\mu \Phi)^\dagger B^{\mu\nu} (D_\nu \Phi) \quad , \tag{3}$$

$$\mathcal{O}_W = \frac{1}{3!} \left(\overrightarrow{W}_\mu^{\ \ \nu} \times \overrightarrow{W}_\nu^{\ \ \lambda} \right) \cdot \overrightarrow{W}_\lambda^{\ \ \mu} \tag{4}$$

induce anomalous triple gauge boson couplings, while[2]

$$\mathcal{O}_{UW} = \frac{1}{v^2} (\Phi^\dagger \Phi - \frac{v^2}{2})\, \overrightarrow{W}^{\mu\nu} \cdot \overrightarrow{W}_{\mu\nu} \quad , \tag{5}$$

$$\mathcal{O}_{UB} = \frac{4}{v^2}\, (\Phi^\dagger \Phi - \frac{v^2}{2}) B^{\mu\nu}\ B_{\mu\nu} \quad , \tag{6}$$

$$\overline{\mathcal{O}}_{UW} = \frac{1}{v^2} (\Phi^\dagger \Phi - \frac{v^2}{2})\, \overrightarrow{W}^{\mu\nu} \cdot \widetilde{\overrightarrow{W}}_{\mu\nu} \quad , \tag{7}$$

$$\overline{\mathcal{O}}_{UB} = \frac{4}{v^2}\, (\Phi^\dagger \Phi - \frac{v^2}{2}) B^{\mu\nu}\ \widetilde{B}_{\mu\nu} \tag{8}$$

create anomalous CP conserving and CP violating Higgs couplings. Note that

$$\widetilde{B}_{\mu\nu} = \frac{1}{2}\, \epsilon_{\mu\nu\rho\sigma} B^{\rho\sigma} \quad , \tag{9}$$

and similarly for $\widetilde{W}_{\mu\nu}$.

Since the operators in (2)-(8) have a dimension higher than four, they will eventually saturate unitarity at sufficiently high energies, unless their locality is tempered by the excitation of new particles. The unitarity constraints for the CP conserving operators shown in (2)-(6) have been derived in [16]. They are given by

$$|f_B| \leq 98 \frac{M_W^2}{s} \qquad , \qquad |f_W| \leq 31 \frac{M_W^2}{s} \quad , \tag{10}$$

$$|\lambda_W| \lesssim 19\, \frac{M_W^2}{s} \quad , \tag{11}$$

[2]In the definition of $\mathcal{O}_{UW}$ and $\mathcal{O}_{UB}$ we have subtracted a trivial contribution to the W and B kinetic energy respectively.

$$|d| \lesssim 17.6\, \frac{M_W^2}{s} + 2.43 \frac{M_W}{\sqrt{s}} \quad , \tag{12}$$

$$-\,236\, \frac{M_W^2}{s} \;+\; 1070\, \frac{M_W^3}{s^{3/2}} \;\lesssim\; d_B \;\lesssim\; 192\, \frac{M_W^2}{s} \;-\; 1123\, \frac{M_W^3}{s^{3/2}} \quad , \tag{13}$$

where s determines the square of the centre of mass energy of the four-boson amplitude where unitarity is first reached.

Here we give the corresponding constraints for the CP violating operators $\overline{\mathcal{O}}_{UW}$ and $\overline{\mathcal{O}}_{UB}$. As in the previous cases, the most important ones arise from the $j = 0$ partial wave amplitudes with vanishing total charge in the s-channel. For the $\overline{\mathcal{O}}_{UB}$ case, the nine channels participating in the transition matrix are ($|\gamma\gamma\pm\pm\rangle$, $|\gamma Z\pm\pm\rangle$, $|ZZ\pm\pm\rangle$, $|ZZLL\rangle$, $|W^-W^+LL\rangle$, $|HH\rangle$), while for the $\overline{\mathcal{O}}_{UW}$ case there are two additional channels given by $|W^-W^+\pm\pm\rangle$. The whole procedure for $\overline{\mathcal{O}}_{UB}$ and $\overline{\mathcal{O}}_{UW}$ is in close analogy to the cases of the operators $\mathcal{O}_{UB}$ and $\mathcal{O}_{UW}$ treated in [16], but this time the tree level amplitudes are complex. For the couplings defined in (1), we find

$$|\overline{d}| \lesssim 18.7\, \frac{M_W^2}{s} + 3.04 \frac{M_W}{\sqrt{s}} \quad , \tag{14}$$

$$|\overline{d}_B| \lesssim 176\, \frac{M_W^2}{s} - 889 \frac{M_W}{\sqrt{s}} \quad . \tag{15}$$

Applying (10)-(15) for $s = 1\,TeV^2$, we get $|f_B| \lesssim 0.6$, $|f_W| \lesssim 0.2$, $|\lambda_W| \lesssim 0.12$, $|d| \lesssim 0.3$, $|d_B| \lesssim 0.7$, $|\overline{d}| \lesssim 0.3$, $|\overline{d}_B| \lesssim 0.7$. There are two remarks to be made concerning these relations. The first is that the constraints for the CP conserving and the CP violating Higgs interactions, derived from ($\mathcal{O}_{UW}$, $\mathcal{O}_{UB}$) and ($\overline{\mathcal{O}}_{UW}$, $\overline{\mathcal{O}}_{UB}$) respectively, are quite similar to each other. The second remark is that the unitarity constraints for the $\overrightarrow{W}_{\mu\nu}$ involving operators ($\mathcal{O}_{W\Phi}$, $\mathcal{O}_W$, $\mathcal{O}_{UW}$, $\overline{\mathcal{O}}_{UW}$), are a factor of 2 to 3 stronger than the corresponding ones for the $B_{\mu\nu}$ involving operators. This means that for similar NP couplings, the new physics forces in the WW channel are considerably stronger than the forces in the ZZ one. A similar situation is known to be valid also for the SM interactions.

To get a feeling of what kind of NP couplings one might expect to appear in (1), we now turn to specific dynamical models. The only requirement in these models is that they always respect $SU(2) \times U(1)$ gauge symmetry and renormalizability. As in the usual SM Lagrangian, no additional discrete symmetry like *e.g.* CP invariance is imposed.

Model A:

In this model, we assume that NP is determined by a complex scalar field Ψ, which has isospin I and hypercharge Y under the $SU(2) \times U(1)$ gauge group. Since Ψ acquires its mass before the electroweak spontaneous breaking, this mass must be large, *i.e.* $M \equiv \Lambda_{NP} \gg v$. The Ψ may also have a hyper-colour $\tilde{N}_c$. The basic renormalizable Lagrangian will then be the sum of the usual SM Lagrangian $\mathcal{L}_{SM}$ and the Lagrangian

$$\mathcal{L}_\psi = D_\mu\Psi^\dagger\, D^\mu\Psi - \Lambda_{NP}^2\Psi^\dagger\Psi + 2g_{\psi 1}(\Psi^\dagger\Psi)(\Phi^\dagger\Phi) + g_{\psi 2}\left[(\Psi^\dagger\widetilde{\Phi})(\widetilde{\Phi}^\dagger\Psi) - (\Psi^\dagger\Phi)(\Phi^\dagger\Psi)\right] \; , \tag{16}$$

describing the interactions of the Ψ field. In writing (16) we have omitted irrelevant $(\Psi^\dagger\Psi)^2$ terms and we have also taken $I \neq 0$ and $(Y \neq 1/6,\ 7/6,\ -5/6\ ,...)$, so that to exclude a direct (Ψ - Φ) mixing and a possible coupling of a single Ψ with either the scalar or the fermion fields of the SM.

The standard techniques may now be used to obtain the effective Lagrangian describing the electroweak interactions at a scale just below Λ_{NP}. This is achieved by integrating out, at the one-loop order, the heavy field Ψ. Thus, by employing the Seeley–DeWitt expansion of the relevant determinant, we obtain the following NP contribution to the electroweak interactions at this scale:

$$
\begin{aligned}
\mathcal{L}_{NP} \;=\; & \frac{(2I+1)\widetilde{N}_c}{(4\pi)^2}\Bigg\{ -\ 2g_{\psi 1}\Lambda^2_{NP}\left(\frac{1}{\epsilon}+1\right)(\Phi^\dagger\Phi)\ +\ \frac{2}{\epsilon}\left(g^2_{\psi 1}+g^2_{\psi 2}\,\frac{I(I+1)}{3}\right)(\Phi^\dagger\Phi)^2 \\
& -\frac{1}{12}\left(\frac{1}{\epsilon}+\frac{g_{\psi 1}v^2}{\Lambda^2_{NP}}\right)\left[\frac{g^2 I(I+1)}{3}\,\overrightarrow{W}_{\mu\nu}\overrightarrow{W}^{\mu\nu}+Y^2 g'^2\ B_{\mu\nu}B^{\mu\nu}\right] \\
& +\frac{8}{6\Lambda^2_{NP}}(g^3_{\psi 1}+g_{\psi 1}g^2_{\psi 2}I(I+1))(\Phi^\dagger\Phi)^3 \\
& +\frac{1}{3\Lambda^2_{NP}}\left(g^2_{\psi 1}+g^2_{\psi 2}\frac{I(I+1)}{3}\right)\partial_\mu(\Phi^\dagger\Phi)\partial^\mu(\Phi^\dagger\Phi) \\
& +\ g^2_{\psi 2}\frac{4I(I+1)}{9\Lambda^2_{NP}}\left[(\Phi^\dagger\Phi)(D_\mu\Phi^\dagger D^\mu\Phi)-(D_\mu\Phi^\dagger\Phi)(\Phi^\dagger D^\mu\Phi)\right] \\
& -\frac{g^2 I(I+1)}{90\Lambda^2_{NP}}\left[g\mathcal{O}_W+\frac{1}{4}\overline{\mathcal{O}}_{DW}+5v^2 g_{\psi 1}\mathcal{O}_{UW}\right] \\
& +\frac{g'Y}{\Lambda^2_{NP}}\left[\frac{2I(I+1)}{9}\,gg_{\psi 2}\mathcal{O}_{BW}-g'g_{\psi 1}Y\frac{v^2}{24}\mathcal{O}_{UB}-\frac{g'Y}{120}\mathcal{O}_{DB}\right]\Bigg\}\ , \qquad (17)
\end{aligned}
$$

where $\epsilon = 2 - n/2$ (with n the number of dimensions) is the usual dimensional regularization parameter, and (4)-(6) are used together with the definitions

$$\overline{\mathcal{O}}_{DW} = 2\,(D_\mu\overrightarrow{W}^{\mu\rho})(D^\nu\overrightarrow{W}_{\nu\rho}) \quad , \qquad (18)$$

$$\mathcal{O}_{DB} = (\partial_\mu B_{\nu\rho})(\partial^\mu B^{\nu\rho}) \quad , \qquad (19)$$

$$\mathcal{O}_{BW} = \frac{1}{2}\,\Phi^\dagger B_{\mu\nu}\overrightarrow{\tau}\cdot\overrightarrow{W}^{\mu\nu}\Phi \quad . \qquad (20)$$

The first three terms in $\mathcal{L}_{NP}$ just renormalize scalar and gauge boson terms already existing in $\mathcal{L}_{SM}$, while the next two indicate an example of how the NP can generate the two unobservable operators $(\Phi^\dagger\Phi)^3$ and $\partial_\mu(\Phi^\dagger\Phi)\partial^\mu(\Phi^\dagger\Phi)$ mentioned in the introduction. The operator $(\Phi^\dagger\Phi)(D_\mu\Phi^\dagger D^\mu\Phi) - (D_\mu\Phi^\dagger\Phi)(\Phi^\dagger D^\mu\Phi)$, as well as the $\overline{\mathcal{O}}_{DW}$, $\mathcal{O}_{DB}$, $\mathcal{O}_{BW}$, given in (18)-(20), should be negligible, according to LEP1 experiments. For reasonable Ψ isospin and hypercharge, this is easily understood for the operators $\overline{\mathcal{O}}_{DW}$ and $\mathcal{O}_{DB}$, whose couplings are proportional to g^2 and g'^2 respectively. The negligible strength of the other two operators just mentioned can be accommodated if we assume that $g_{\psi 2}$ (defined in (16)) is negligible.

The interesting thing about Model A is that there is nothing that prohibits $g_{\psi 1}$ (also defined in (16)) to be large. And if this does happen, then the operators $\mathcal{O}_{UW}$ and $\mathcal{O}_{UB}$ will be proportionally enhanced by NP. Unfortunately, an analogous enhancement for $\mathcal{O}_W$ is not so easy. The $\mathcal{O}_W$ coupling expected perturbatively satisfies $\lambda_W \sim g^2$, and it should therefore be similar to the coupling of the strongly constrained operator $\overline{\mathcal{O}}_{DW}$ (see (18)). Only if $\mathcal{O}_W$ is somehow non-perturbatively enhanced with respect to $\overline{\mathcal{O}}_{DW}$, by a mechanism like the one discussed in [17], we could hope that it would become observable.

Therefore, out of the seven operators appearing in (1), Model A favours only $\mathcal{O}_{UW}$ and $\mathcal{O}_{UB}$, and to a lesser extent $\mathcal{O}_W$. The couplings of the Higgs involving operators $\mathcal{O}_{UW}$ and $\mathcal{O}_{UB}$ depend on the unknown physics of the scalar sector. Thus, these two later operators really teach us something about the mechanism that breaks spontaneously the gauge symmetry. On the contrary the purely gauge dependent operator $\mathcal{O}_W$ seems naturally suppressed at the perturbative level, by the small coupling g. Nevertheless, it is at least generated in this model. Note that if Ψ had $Y = 0$, then only the custodially $SU(2)_c$ invariant operators $\mathcal{O}_{UW}$ and $\mathcal{O}_W$ would have appeared[3] [17].

Model B:

We now turn to Model B where NP is determined instead by a fermion field F whose left and right component have the same isospin I and hypercharge Y. Because of the vectorial character of the model, there are no anomalies and F acquires its mass before the spontaneous electroweak breaking takes place. Hence, we can assume that F has a very large mass Λ_{NP} and possibly also a hyper-colour $\widetilde{N}_c$. To construct the basic renormalizable Lagrangian we should now add to $\mathcal{L}_{SM}$ the term

$$\mathcal{L}_F = i\overline{F}(\partial\!\!\!/ + ig\overrightarrow{W\!\!\!\!/}\cdot\overrightarrow{t} + ig'YB\!\!\!/)F - \Lambda_{NP}\overline{F}F \quad , \tag{21}$$

with $\overrightarrow{t}$ denoting the isospin I matrices. In writing (21) we have excluded a discrete set of hypercharge and isospin values which would allow a coupling of F with the SM fermions and possibly also with the standard Higgs.

Integrating the fermion loop as before [18], we get at the scale Λ_{NP}:

$$\begin{aligned}\mathcal{L}_{NP} \;=\; \frac{(2I+1)\widetilde{N}_c}{(4\pi)^2}\Bigg\{ &-\frac{g^2 I(I+1)}{9\epsilon}\,\overrightarrow{W}_{\mu\nu}\overrightarrow{W}^{\mu\nu} \;-\; \frac{Y^2 g'^2}{3\epsilon}\,B_{\mu\nu}B^{\mu\nu} \\ &+\frac{g^2 I(I+1)}{45\Lambda_{NP}^2}\Big[g\mathcal{O}_W - \overline{\mathcal{O}}_{DW}\Big] \;-\; \frac{g'^2Y^2}{15\Lambda_{NP}^2}\,\mathcal{O}_{DB}\Bigg\} \;.\end{aligned} \tag{22}$$

The only interesting operator generated in this case is $\mathcal{O}_W$, about which though (as well as about the unwanted operators $\overline{\mathcal{O}}_{DW}$ and $\mathcal{O}_{DB}$), the same remarks apply as in Model A. It seems that if NP only includes fermionic new degrees of freedom, we cannot learn much about the scalar sector by studying the anomalous bosonic couplings. To reiterate on this we thus turn to Model B′.

[3] Provided of course that we still keep the assumption $g_{\psi 2} \sim 0$.

Model B′ :

To the preceding NP spectrum we just add a heavy ($M_s \gg v$) scalar field S^0 with vanishing isospin and hypercharge. Then the most general renormalizable interaction to be added to $\mathcal{L}_{SM}$ becomes

$$\begin{aligned}\mathcal{L}_{FS} &= i\overline{F}\not{D}F - \Lambda_{NP}\overline{F}F + g_f S^0\overline{F}F + ig_{f'}S^0\overline{F}\gamma_5 F \\ & g_\phi M_s S^0\Phi^\dagger\Phi + \frac{1}{2}(\partial S)^2 - \frac{M_s^2}{2}(S^0)^2 \quad ,\end{aligned} \tag{23}$$

where irrelevant $(S^0)^3$ and $(S^0)^4$ terms have been omitted. Integrating out first the heavy F field, and then substituting S^0 to $g_\phi\Phi^\dagger\Phi/M_S$, we find that NP generates, in addition to the terms appearing in (22), contributions also from all four purely Higgs operators shown in (1), *i.e.* $\mathcal{O}_{UW}$, $\mathcal{O}_{UB}$, $\overline{\mathcal{O}}_{UW}$, $\overline{\mathcal{O}}_{UB}$. Restricting for simplicity to $I = 1/2$ for the isospin of the F fermion, the couplings of the CP conserving operators are expressed as

$$d = -\left(\frac{g^2 v^2 \widetilde{N}_c}{48\pi^2\Lambda_{NP}M_s}\right) g_f g_\phi \quad , \tag{24}$$

$$d_B = -\left(\frac{g'^2 v^2 Y^2 \widetilde{N}_c}{48\pi^2\Lambda_{NP}M_s}\right) g_f g_\phi \quad , \tag{25}$$

while those of the CP violating ones satisfy

$$\frac{\overline{d}}{d} = \frac{\overline{d}_B}{d_B} = -\frac{g_{f'}}{g_f} \quad . \tag{26}$$

Note that for vanishing hypercharge Y for the F fermion, only the custodially $SU(2)_c$ invariant operators O_{UW} and $\overline{O}_{UW}$ would be generated [17]. Since the g_f and $g_{f'}$ couplings in (23) are a priori on the same footing, we conclude that in this model all four operators $\mathcal{O}_{UW}$, $\mathcal{O}_{UB}$, $\overline{\mathcal{O}}_{UW}$, $\overline{\mathcal{O}}_{UB}$ can be generated with appreciable couplings. We also note that if the scalar boson S were chosen instead to be isovector, then the model would have generated the operator $\mathcal{O}_{BW}$ and its CP violating analogue. Since the last two operators are very strongly constrained from existing experiments, we would conclude that such a situation is disfavoured.

The above considerations lead to the conclusion that the question whether one of the operators ($\mathcal{O}_{UW}$, $\mathcal{O}_{UB}$, $\overline{\mathcal{O}}_{UW}$, $\overline{\mathcal{O}}_{UB}$) will be generated or not is intimately connected with the nature of the mechanism responsible for the spontaneous breaking of the gauge symmetry. The experimental search for such an operator will teach us something on how the spontaneous symmetry breaking works. Our models imply also that, to a lesser extent, $\mathcal{O}_W$ can also be generated by NP; but this operator seems to be rather independent of the scalar sector. Finally the other two operators in (1), namely $\mathcal{O}_{W\Phi}$ and $\mathcal{O}_{B\Phi}$, were never generated in our models.

It is unnecessary to state that we take these models only as indicative. The laws of New Physics are certainly much more elaborate than our toy models suggest. It could also be that the residual interactions below the NP scale Λ_{NP} not only involve weak bosons but also heavy quarks, *i.e.* the third family. In [3] we have established the list of 28 $dim = 6$ operators involving the third family, 14 of them involving the t_R field (which in SM is associated to the top mass generation) and we showed that some of them could also be at the origin of the departure of $Zb\bar{b}$ from SM predictions.

In any case an active experimental search at LEP2 and at higher energy colliders should be made in order to identify any of the operators we have discussed. Once any of them is found, then (as in the old Fermi theory), the unitarity constraints presented above may help deciding how far we are from the energy region where some new degrees of freedom should start being excited.

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References

[1] See *e.g.* the talk given by G. Altarelli at the Rome Conference on "Phenomenology of Unification from Present to Future", Proc. of the EPS Conference on High Energy Physics, Marseille, France, 1993 CERN-TH.7319/94, CERN-TH.7045/93 July 1993. G. Quast, *ibid.*; J. Lefrançois *ibid.*; D. Schaile, CERN-PPE/93-213 (1993); LEP Electroweak Working Group, CERN-PPE/93-157 (1993); The LEP Collaborations, ALEPH, DELPHI, L3, OPAL and the LEP Electroweak Working Group, CERN/PPE/94-187 ; G. Altarelli, CERN-TH.7464/94.

[2] D. Schaile, presented at the 27th Int. Conf. on High Energy Phys., Glasgow (1994). J. Erler and P. Langacker, Univ of Pennsylvania preprint UPR-0632T (1994).

[3] G.J. Gounaris, F.M. Renard and C. Verzegnassi, Preprint PM/94-44 and THES-TP 95/01, hep-ph/9501362.

[4] M.S. Chanowitz, Annu. Rev. Nucl. Part. Sci. **38** (1988) 323.

[5] W. Buchmüller and D. Wyler, Nucl. Phys. **B268** (1986) 621; C.J.C. Burgess and H.J. Schnitzer, Nucl. Phys. **B228** (1983) 454; C.N. Leung, S.T. Love and S. Rao Z. Phys. **C31** (1986) 433.

[6] R. Casalbuoni, S. De Curtis, D. Dominici and R. Gatto, Phys. Lett. **B155** (1985) 95 and Nucl. Phys. **B282** (1987) 235. J.L. Kneur and D. Schildknecht, Nucl. Phys. **B357** (357) 1991; J.L. Kneur, M. Kuroda and D. Schildknecht, Phys. Lett. **B262** (93) 1991 .

[7] J. Layssac, F.M. Renard and C. Verzegnassi, Phys. Lett. **B287** (1992) 267 and Phys. Rev. **D49** (1994) 3650.

[8] K. Hagiwara, S. Ishihara, R. Szalapski and D. Zeppenfeld, Phys. Lett. **B283** (1992) 353 and Phys. Rev. **D48** (1993) 2182.

[9] K. Hagiwara and M.L. Stong, Z. Phys. **C62** (1994) 99.

[10] A. De Rújula, M.B. Gavela, P. Hernandez and E. Masso, Nucl. Phys. **B384** (1992) 3.

[11] A. Bilal, E. Massó and A. De Rújula, Nucl. Phys. **B355** (1991) 549; A. De Rújula, M. Gavela, O. Pène and F.J. Vegas, Nucl. Phys. **B357** (1991) 357.

[12] M. Bilenky, J.L. Kneur, F.M. Renard and D. Schildknecht, Nucl. Phys. **B409** (1993) 22 and **B419** (1994) 240.

[13] G.J Gounaris, J. Layssac, G. Moultaka and F.M. Renard, Int. J. Mod. Phys. **A8** (1993) 3285.

[14] M. Krämer, J. Kühn, M.L. Stong and P.M. Zerwas, Z. Phys. **C64** (1994) 21.

[15] G.J Gounaris et al., in preparation.

[16] G.J. Gounaris, J. Layssac and F.M. Renard, Phys. Lett. **B332** (1994) 146 ; G.J. Gounaris, J. Layssac, J.E. Paschalis and F.M. Renard, Montpellier preprint PM/94-28, to appear in Z. Phys.

[17] G.J. Gounaris, F.M. Renard and G. Tsirigoti, Phys. Lett. **B338** (1994) 51 .

[18] R.D. Ball, Phys. Re. **C182** (1989) 1.