We study the indirect effects of new physics on the phenomenology of the recently discovered "Higgs-like" particle. In a model independent framework these effects can be parametrized in terms of an effective Lagrangian at the electroweak scale. In a theory in which the $SU(2)_L \times U(1)_Y$ gauge symmetry is linearly realized they appear at lowest order as dimension-six operators, containing all the SM fields including the light scalar doublet, with unknown coefficients. We discuss the choice of operator basis which allows us to make better use of all the available data on the new state, triple gauge boson vertex and electroweak precision tests, to determine the coefficients of the new operators. We illustrate our present knowledge of these by performing a global fit to the existing data which allows simultaneous determination of the eight relevant parameters quantifying the Higgs couplings to gluons, electroweak gauge bosons, bottom quarks, and tau leptons. We find that for all scenarios considered the standard model predictions for each individual Higgs coupling and observable are within the corresponding 68% CL allowed range. We finish by commenting on the implications of the results for unitarity of processes at higher energies.

Note added: The analysis has been updated with all the public data available by October 2013. Updates of this analysis are provided at the web site http://hep.if.usp.br/Higgs as well as new versions of this manuscript.
of the low–energy theory.

With the present data we can proceed by assuming that the observed state belongs indeed to a light electroweak doublet scalar and that the $SU(2)_L \otimes U(1)_Y$ symmetry is linearly realized in the effective theory $^{12,13}$. Barring effects associated with violation of total lepton number, the lowest order operators which can be built are of dimension six. The coefficients of these dimension–six operators parametrize our ignorance of the new physics effects in the Higgs phenomenology and our task at hand is to determine them using the available data. This bottom–up approach has the advantage of minimizing the amount of theoretical hypothesis when studying the Higgs couplings.

Following this approach we start by listing in Sec. II the most general set of dimension–six operators which involve triple couplings of the low energy scalar to the SM gauge bosons and fermions and can affect the present Higgs data. The list is redundant and, at any order, the operators listed are related by the equations of motion (EOM). This allows for a freedom of choice in the election of the basis of operators to be used in the analysis. We will argue in Sec. III that in the absence of any “a priori” illumination on the form of the new physics the most sensible choice of basis should contain operators whose coefficients are more easily related to existing data from other well tested sectors of the theory. This will reduce to eight the number of operators testable with an analysis of the existing Higgs data. We proceed then to briefly describe in Sec. III the technical details of such analysis. The status of this exercise with the most up–to–date experimental results is presented in Sec. IV which updates the analysis in Ref. [9] also extending the previous analysis by including modifications of the Higgs couplings to fermions. We summarize our conclusions in Sec. V.

II. EFFECTIVE LAGRANGIAN FOR HIGGS INTERACTIONS

In order to probe the Higgs couplings we parametrize the deviations from the SM predictions in terms of effective Lagrangians. Here, we assume that the low energy theory exhibits all the symmetries of the SM and that it contains only the SM degrees of freedom. Furthermore, we consider that the recently observed state belongs to an $SU(2)_L$ doublet $^1$. We further assume that the present precision of the data allows us to parametrize the deviations from the SM predictions by operators of dimension up to six, i.e.

$$\mathcal{L}_{\text{eff}} = \sum_n \frac{f_n}{\Lambda^2} \mathcal{O}_n ,$$

(1)

where the dimension–six operators $\mathcal{O}_n$ involve gauge–boson, Higgs–boson and/or fermionic fields with couplings $f_n$ and where $\Lambda$ is a characteristic scale. Moreover, we assumed the $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ SM local symmetry as well as the $\mathcal{O}_n$ operators to be $P$ and $C$ even and the conservation of baryon and lepton numbers.

Our first task is to fix the basis of dimension–six operators that is suitable to study the Higgs couplings. Of all dimension–six operators just 59 of them, up to flavor and Hermitian conjugation, are enough to generate the most general $S$–matrix elements consistent with baryon number conservation and the SM gauge symmetry $^{20}$. Before deciding the operators used in our analyses, let us discuss the dimension–six effective interactions that modify the Higgs coupling to gauge bosons and to fermions.

A. Higgs interactions with gauge bosons

To start, we consider the following eight $P$ and $C$ even dimension–six operators that modify the Higgs couplings to the electroweak gauge bosons, and one operator containing gluons $^{12,13}$:

$$\mathcal{O}_{GG} = \Phi \Phi G_{\mu \nu}^a G^{a \mu \nu} , \quad \mathcal{O}_{WW} = \Phi \Phi W_\mu W^{\mu \nu} \Phi , \quad \mathcal{O}_{BB} = \Phi \Phi \tilde{B}_\mu \tilde{B}^{\mu \nu} \Phi ,$$

$$\mathcal{O}_{BW} = \Phi \Phi B_\mu \tilde{W}^{\mu \nu} \Phi , \quad \mathcal{O}_{W} = (D_\mu \Phi) \tilde{W}^{\mu \nu} (D_\nu \Phi) , \quad \mathcal{O}_{B} = (D_\mu \Phi) \tilde{B}^{\mu \nu} (D_\nu \Phi) ,$$

$$\mathcal{O}_{\Phi,1} = (D_\mu \Phi) \Phi \Phi (D^\mu \Phi) , \quad \mathcal{O}_{\Phi,2} = \frac{i}{2} \partial^{\mu} \left( \Phi \Phi \right) \partial_{\mu} \left( \Phi \Phi \right) , \quad \mathcal{O}_{\Phi,4} = (D_\mu \Phi) \Phi (D^\mu \Phi) ,$$

(2)

where we denoted the Higgs doublet by $\Phi$ and its covariant derivative is $D_\mu \Phi = (\partial_\mu + i \frac{1}{2} g' B_\mu + ig \frac{g'}{2} W_\mu^a) \Phi$ in our conventions. The hatted field strengths are defined as $\hat{B}_{\mu \nu} = i \frac{g}{2} B_{\mu \nu}$ and $\hat{W}_{\mu \nu} = i \frac{g'}{2} \sigma^a W_{\mu \nu}^a$. Moreover, we denote the

---

$^1$ This implies that the new physics decouples when the cut–off $\Lambda \to \infty$.
SU(2)_L (U(1)_Y) gauge coupling as $g (g')$ and the Pauli matrices as $\sigma^\alpha$. Our conventions are such that

$$W^\pm_\mu = \frac{1}{\sqrt{2}} (W^1_\mu + i W^2_\mu), \quad Z^\text{SM}_\mu = \frac{1}{\sqrt{g^2 + g'^2}} (g W^3_\mu - g' B_\mu), \quad A^\text{SM}_\mu = \frac{1}{\sqrt{g^2 + g'^2}} (g' W^3_\mu + g B_\mu).$$

In the unitary gauge the Higgs field is written as

$$\Phi = \frac{1}{\sqrt{2}} \begin{pmatrix} v + h(x) \end{pmatrix},$$

where $v$ is its vacuum expectation value (vev).

For the sake of completeness of our discussion, it is interesting to introduce the operator that contains exclusively Higgs fields

$$O_{\Phi,3} = \frac{1}{3} (\Phi^\dagger \Phi)^3$$

that gives an additional contribution to the Higgs potential

$$\mu_0^2 (\Phi^\dagger \Phi) + \lambda_0 (\Phi^\dagger \Phi)^2 - \frac{f_{\Phi,3}}{3\Lambda^2} (\Phi^\dagger \Phi)^3.$$

This effective operator leads to a shift of the minimum of the Higgs potential with respect to the SM result

$$v^2 = -\frac{\mu_0^2}{\lambda_0} \left( 1 + \frac{v^2}{4\Lambda^2} \frac{f_{\Phi,3}}{\lambda_0} \right) = v_0^2 \left( 1 + \frac{v^2}{4\Lambda^2} \frac{f_{\Phi,3}}{\lambda_0} \right) \tag{7}$$

The operators $O_{\Phi,1}, O_{\Phi,2},$ and $O_{\Phi,4}$ contribute to the kinetic energy of the Higgs boson field $h$ so we need to introduce a finite wave function renormalization in order to bring the Higgs kinetic term to the canonical form

$$H = h \left[ 1 + \frac{v^2}{2\Lambda^2} (f_{\Phi,1} + 2f_{\Phi,2} + f_{\Phi,4}) \right]^{1/2}. \tag{8}$$

Furthermore, the operators $O_{\Phi,j}$ ($j = 1, 2, 3, 4$) also alter the Higgs mass according to

$$M_H^2 = 2\lambda_0 v^2 \left[ 1 - \frac{v^2}{2\Lambda^2} \left( f_{\Phi,1} + 2f_{\Phi,2} + f_{\Phi,4} \right) \right],$$

where we have expanded to linear order in the $f_i$ coefficients.

The operator $O_{BW}$ contributes at tree level to $Z\gamma$ mixing, therefore, the mass eigenstates are

$$Z_\mu = \left[ 1 - \frac{g^2 g'^2}{2(g^2 + g'^2) \Lambda^2 f_{BW}} \right]^{-1/2} Z^\text{SM}_\mu, \tag{9}$$

$$A_\mu = \left[ 1 + \frac{g^2 g'^2}{2(g^2 + g'^2) \Lambda^2 f_{BW}} \right]^{-1/2} A^\text{SM}_\mu - \left[ \frac{gg' (g^2 - g'^2)}{4(g^2 + g'^2) \Lambda^2 f_{BW}} \right] Z^\text{SM}_\mu. \tag{10}$$

The operators $O_{BW}, O_{\Phi,1}, O_{\Phi,3}$ and $O_{\Phi,4}$ also have an impact on the electroweak gauge boson masses. Expanding to linear order in the $f_i$ coefficients they read:

$$M_Z^2 = \frac{g^2 + g'^2}{4} v^2 \left[ 1 + \frac{v^2}{2\Lambda^2} (f_{\Phi,1} + f_{\Phi,4} - \frac{g^2 g'^2}{(g^2 + g'^2) f_{BW}}) \right], \tag{11}$$

$$M_W^2 = \frac{g^2}{4} v^2 \left[ 1 + \frac{v^2}{2\Lambda^2} f_{\Phi,4} \right]. \tag{12}$$

Notice that $O_{BW}$ and $O_{\Phi,1}$ contribute to the $Z$ mass but not to the $W$ mass, therefore, violating the custodial $SU(2)$ symmetry and contributing to $T$ (or $\Delta \rho$).

In our calculations we will always use as inputs the measured values of $G_F, M_Z$ and $\alpha$, where the electromagnetic coupling is evaluated at zero momentum. Furthermore, when convenient, we will also absorb part of the tree–level
renormalization factors by using the measured value of $M_W$. In particular using $\frac{G_F}{\sqrt{2}} = \frac{g^2}{8M_W}$ and Eqs. (11) and (12) we obtain that

$$v = \left(\sqrt{2}G_F\right)^{-1/2} \left(1 - \frac{v^2}{4\Lambda^2}f_{\phi,4}\right)$$

$$M_Z^2 = \left(\sqrt{2}G_F\right)^{-1} \frac{g^2}{4c^2} \left(1 + \frac{v^2}{2\Lambda^2}f_{\phi,1} - \frac{g^2g^2}{2(g^2 + g'^2)}\frac{v^2}{\Lambda^2}f_{BW}\right)$$

where we have denoted by $c \equiv g/\sqrt{g^2 + g'^2}$ the tree level cosine of the SM weak mixing angle.

The dimension–six effective operators in Eq. (2) give rise to Higgs interactions with SM gauge–boson pairs that take the following form in the unitary gauge.

$$\mathcal{L}_{eff}^{HV} = g_{Hgg} H G_{\mu\nu} G^{\mu\nu} + g_{H\gamma\gamma} H A_{\mu\nu} A^{\mu\nu} + g_{HZZ} Z_{\mu\nu} Z^{\mu\nu} + g_{H\gamma\gamma} H A_{\mu\nu} Z^{\mu\nu}$$

$$+ g_{HZZ} Z_{\mu\nu} Z^{\mu\nu} + g_{HWW} W_+^{\mu\nu} W_-^{\mu\nu} + g_{HWW} H W_+^{\mu\nu} W_-^{\mu\nu}$$

$$+ g_{HWW} W_+^{\mu\nu} W_-^{\mu\nu} + g_{HWW} H W_+^{\mu\nu} W_-^{\mu\nu}$$

where $V_{\mu\nu} = \partial_\mu V_\nu - \partial_\nu V_\mu$ with $V = A, Z, W,$ and $G$. The effective couplings $g_{Hgg}, g_{H\gamma\gamma}, g_{HZZ}, g_{HWW}$ and $g_{HZZ}$ are related to the coefficients of the operators appearing in Eq. (1) through

$$g_{Hgg} = \frac{f_{GG}^G}{2\Lambda^2} = \frac{\alpha_s f_g^v}{8\pi \Lambda^2},$$

$$g_{H\gamma\gamma} = \frac{(g^2 + g'^2)^2}{2\Lambda^2} \frac{f_{BB} + f_{BW} - f_{BW}}{2},$$

$$g_{HZZ}^{(1)} = \frac{g^2}{2\Lambda^2} \frac{s(f_{BW} - f_B)}{2c},$$

$$g_{HZZ}^{(2)} = \frac{g^2}{2\Lambda^2} \frac{f_{BB} + f_{BW} - f_{BW}}{2c},$$

$$g_{HZZ}^{(3)} = \frac{g^2}{2\Lambda^2} \frac{c^2 f_{BW} + s^2 f_{BW}}{2c^2},$$

$$g_{HWW}^{(1)} = \frac{g^2}{2\Lambda^2} \frac{f_W}{2},$$

$$g_{HWW}^{(2)} = \frac{g^2}{2\Lambda^2} \frac{f_{WW}}{2},$$

$$g_{HWW}^{(3)} = \frac{g^2}{2\Lambda^2} \frac{c^2 f_{BW} + s^2 f_{BW}}{2c^2},$$

where $s \equiv g'/\sqrt{g^2 + g'^2}$ stands for the tree level sine of the SM weak mixing angle. For convenience, we rescaled the coefficient $f_{GG}$ of the gluon–gluon operator by a loop factor $-\alpha_s/(8\pi)$ such that an anomalous gluon–gluon coupling $f_g \sim O(1 - 10)$ gives a contribution comparable to the SM top loop. Furthermore, we have kept the normalization commonly used in the pre-LHC studies for the operators involving electroweak gauge bosons. Notice that the general expressions above reproduce in the different cases considered those of Refs. [16, 18, 21].
B. Higgs interactions with fermions

The dimension–six operators modifying the Higgs interactions with fermion pairs are [20]

\[
\begin{align*}
O_{e\Phi,ij} &= (\Phi^\dagger \Phi) (\bar{L}_i \Phi e_{Rj}), \\
O_{\Phi_{L,ij}} &= \Phi^\dagger (iD^\mu_{\nu} \Phi) (\bar{L}_i \sigma_\mu \sigma_\nu L_j), \\
O_{u\Phi,ij} &= (\Phi^\dagger \Phi) (\bar{Q}_i \Phi u_{Rj}), \\
O_{Q_{L,ij}} &= \Phi^\dagger (iD^\mu_{\nu} \Phi) (\bar{Q}_i \gamma_\mu \nu Q_j), \\
O_{d\Phi,ij} &= (\Phi^\dagger \Phi) (\bar{Q}_i \Phi d_{Rj}), \\
O_{Q_{d,ij}} &= \Phi^\dagger (iD^\mu_{\nu} \Phi) (\bar{d}_i \gamma_\mu \nu d_{Rj}),
\end{align*}
\]

where we define \( \tilde{\Phi} = \sigma_2 \Phi^*, \) \( \Phi^\dagger D^\mu_{\nu} \Phi = \Phi^\dagger D^\mu_{\nu} \Phi - (D^\mu \Phi)^\dagger \Phi \) and \( \Phi^\dagger D^\mu_{\nu} \Phi = \Phi^\dagger \sigma^a D^\mu \Phi - (D^\mu \Phi)^\dagger \sigma^a \Phi. \) We use the notation of \( L \) for the lepton doublet, \( Q \) for the quark doublet and \( f_R \) for the \( SU(2) \) singlet fermions, where \( i, j \) are family indices. Notice that, unlike the Higgs–gauge boson operators of the previous subsection, not all Higgs–fermion operators listed above are Hermitian.

In Eq. (17) we have classified the operators according to the number of Higgs fields they contain. In a first set, \( O_{\Phi_{L,ij}} \), the operators exhibit three Higgs fields and after spontaneous symmetry breaking they lead to modifications of the SM Higgs Yukawa couplings. The second set, \( O_{\Phi_{Q,ij}}^{(1)} \), contains operators presenting two Higgs fields and one covariant derivative, and consequently, they contribute to Higgs couplings to fermion pairs which also modify the charged weak interactions. The third set, \( O_{\Phi_{d,ij}}^{(3)} \), similar to the second, also leads to modifications of the fermionic neutral and charged current interactions.

Operators \( O_{f\Phi,ij} \) renormalize fermion masses and mixing, as well as modify the Yukawa interactions. In the SM, these interactions take the form

\[
\mathcal{L}_{Yuk} = -y^d_{ij} \bar{L}_i \Phi e_{Rj} - y^u_{ij} \bar{Q}_i \Phi d_{Rj} - y^e_{ij} \bar{Q}_i \Phi u_{Rj} + \text{h.c.},
\]

while the dimension–six modifications of the Yukawa interactions are

\[
\mathcal{L}_{\text{eff}}^{H} = \frac{f_{\Phi} f_{\Phi}^{\dagger}}{\Lambda^2} O_{\Phi_{L,ij}} + \frac{f_{\Phi} f_{\Phi}^{\dagger}}{\Lambda^2} O_{\Phi_{Q,ij}} + \frac{f_{\Phi} f_{\Phi}^{\dagger}}{\Lambda^2} O_{\Phi_{d,ij}} + \text{h.c.}
\]

where a sum over the three families \( i, j = 1, 2, 3 \) is understood. After spontaneous symmetry breaking and prior to the finite Higgs wave function renormalization in Eq. (17), Eqs. (18) and (19) can be conveniently decomposed in two pieces \( \mathcal{L}_0 \) and \( \mathcal{L}_1 \) given by

\[
\mathcal{L}_0 = \frac{1}{\sqrt{2}} \bar{d}_L \left[ -y^d + \frac{v^2}{2\Lambda^2} f_{\Phi} \right] d_R (v+h) + \frac{1}{\sqrt{2}} \bar{u}_L \left[ -y^u + \frac{v^2}{2\Lambda^2} f_{\Phi} \right] u_R (v+h) + \frac{1}{\sqrt{2}} \bar{e}_L \left[ -y^e + \frac{v^2}{2\Lambda^2} f_{\Phi} \right] e_R (v+h) + \text{h.c.},
\]

and

\[
\mathcal{L}_1 = \frac{1}{\sqrt{2}} \bar{d}_L f_{\Phi} d_R h + \frac{1}{\sqrt{2}} \bar{u}_L f_{\Phi} u_R h + \frac{1}{\sqrt{2}} \bar{e}_L f_{\Phi} e_R h + \text{h.c.},
\]

where \( f_{L,R} = (f_{L,R1}, f_{L,R2}, f_{L,R3})^T \) with \( f = u, \) or \( d \) and \( e \) and \( y^d \) and \( y^f \) are 3×3 matrices in generation space.

\( \mathcal{L}_0 \) is proportional to the mass term for the fermions and in the mass basis leads to the SM–like Higgs–fermion interactions with renormalized fermion masses and quark weak mixing. On the other hand, generically, the new interactions contained in \( \mathcal{L}_1 \) are not necessarily flavor diagonal in the mass basis unless \( f_{\Phi} \propto y^f. \)

Altogether the \( H \bar{f} f \) couplings in the fermion mass basis and after renormalization of the Higgs wave function in Eq. (18) can be written as

\[
\mathcal{L}^{H \bar{f} f} = g_{fij} f_{L,R} f_{L,R} + \text{h.c.},
\]

2. Since we are not adding right-handed neutrinos to the fermion basis or allowing for \( L \) violating dimension–five operators, the couplings to the charged leptons can be chosen to be generation diagonal in the mass basis as in the SM.
with

\[ g_{Hij}^f = -\frac{m_f}{v} \delta_{ij} \left[ 1 - \frac{v^2}{4\Lambda^2} (f_{\Phi,1} + 2f_{\Phi,2} + f_{\Phi,4}) \right] + \frac{v^2}{\sqrt{2}\Lambda^2} f'_{\Phi,ij} \]  

(23)

where we denoted the physical masses by \( m_f \) and \( f'_{\Phi,ij} \) are the coefficients of the corresponding operators in the mass basis. In what follows we will denote all these coefficients without the prime.

C. The right of choice

In the effective Lagrangian framework not all operators at a given order are independent as they can be related by the use of the classical equations of motion (EOM) of the SM fields. The invariance of the physical observables under the associated operator redefinitions is guaranteed as it has been proved that operators connected by the EOM lead to the same \( S \)-matrix elements [22]. In a top–bottom approach, when starting from the full theory and integrating out heavy degrees of freedom to match the coefficients of the higher dimension operators at low energies it is convenient not to choose a minimal set of operators in order to guarantee that the operators generated by the underlying theory can be easily identified [23]. However, in a bottom–up approach when we use the effective Lagrangians to obtain bounds on generic extensions of the SM, we must choose a minimum operator basis to avoid parameters combinations that can not be probed.

In our case at hand, we have to take into account the SM EOM which imply that not all the operators in Eqs. (2) and (17) are independent. In particular the EOM for the Higgs field and the electroweak gauge bosons lead to three relations between the operators:

\[ 2O_{\Phi,2} + 2O_{\Phi,4} = \sum_{ij} \left( y_{ij}^c (O_{\Phi,ij}) + y_{ij}^d (O_{\Phi,ij})^\dagger + h.c. \right) - \frac{\partial V(h)}{\partial h} , \]  

(24)

\[ 2O_B + O_{BW} + O_{BB} + g^2 \left( O_{\Phi,1} - \frac{1}{2} O_{\Phi,2} \right) = -\frac{g^2}{2} \sum_i \left( \frac{1}{2} O^{(1)}_{\Phi,L,ii} + \frac{1}{6} O^{(1)}_{\Phi,ii} - O^{(1)}_{\Phi_e,ii} + \frac{2}{3} O^{(1)}_{\Phi_u,ii} - \frac{1}{3} O^{(1)}_{\Phi_d,ii} \right) , \]  

(25)

\[ 2O_W + O_{BW} + O_{WW} + g^2 \left( O_{\Phi,1} - \frac{1}{2} O_{\Phi,2} \right) = -\frac{g^2}{4} \sum_i \left( O^{(3)}_{\Phi,L,ii} + O^{(3)}_{\Phi_Q,ii} \right) . \]  

(26)

These constraints allow for the elimination of three operators listed in Eqs. (2) and (17).

At this point we are faced with the decision of which operators to leave in the basis to be used in the analysis of the Higgs data; different approaches can be followed in doing so. Again, in a top–bottom approach in which some \( a \) \( priori \) knowledge is assumed about the beyond the SM theory one can use this theoretical prejudice to choose the basis. For example if the UV completion of the SM is a given gauge theory, it is possible to predict whether a given operator is generated at tree level or at loop level [24]. One may then be tempted to keep those in the basis as basis. For example if the UV completion of the SM is a given gauge theory, it is possible to predict whether a given operator is generated at tree level or at loop level [24]. One may then be tempted to keep those in the basis as

\[ \mathcal{L}_{WWV} = -ig_{WWV} \left\{ g_1^V \left( W^\nu_{\mu\nu} W^{-\mu} - W^\nu_{\mu\nu} W^{-\mu} \right) + \kappa_V W^+_\nu V^\mu + \frac{\lambda_V}{m_W^2} W^+_{\mu\nu} W^{-\nu\mu} V^\mu \right\} , \]  

(27)
where $g_{WW\gamma} = e = g_s$, $g_{WWZ} = g_c$. In general these vertices involve six dimensionless couplings $g_1^{Y}, \kappa_V$, and $\lambda_V$ ($V = \gamma$ or $Z$). Notwithstanding the electromagnetic gauge invariance requires that $g_1^{Z} = 1$, while the three remaining couplings are related to the dimension-six operators $\mathcal{O}_B, \mathcal{O}_W, \mathcal{O}_{BW}$, and $\mathcal{O}_{\Phi,1}$

$$
\Delta g_1^Z = g_1^Z - 1 = \frac{g^2 v^2}{8 c^2 \Lambda^2} \left( f_W + 2 \frac{s^2}{c^2 - s^2} f_{BW} \right) - \frac{1}{4(c^2 - s^2)} f_{\Phi,1} \frac{v^2}{\Lambda^2}, \\
\Delta \kappa_\gamma = \kappa_\gamma - 1 = \frac{g^2 v^2}{8 \Lambda^2} \left( f_W + f_B - 2 f_{BW} \right), \\
\Delta \kappa_Z = \kappa_Z - 1 = \frac{g^2 v^2}{8 c^2 \Lambda^2} \left( c^2 f_W - s^2 f_B + \frac{4 s^2 c^2}{c^2 - s^2} f_{BW} \right) - \frac{1}{4(c^2 - s^2)} f_{\Phi,1} \frac{v^2}{\Lambda^2},
$$

(28)

Next we notice that the $Z$ and $W$ couplings to fermions as well as the oblique parameters $S, T, U$ are in agreement with the SM at the per mil to per cent level [28]. These results impose severe constraints on the operators which modify these observables: $\mathcal{O}_{\Phi,1}^{(1)}, \mathcal{O}_{\Phi,1}^{(3)}, \mathcal{O}_{BW}$ and $\mathcal{O}_{\Phi,1}$. For example $\mathcal{O}_{BW}$ and $\mathcal{O}_{\Phi,1}$ give a tree level contribution to $S$ and $T$ [14, 15, 25, 28]

$$
\alpha \Delta S = c^2 \frac{v^2}{\Lambda^2} f_{BW} \\ \alpha \Delta T = \frac{1}{2} \frac{v^2}{\Lambda^2} f_{\Phi,1}.
$$

(29)

However, in order to take full advantage of these electroweak precision observables (EWPO) we must be sure that there is no combination of the anomalous operators whose contribution at the tree level to EWPO cancels out, i.e. we must avoid the existence of blind directions [14, 30].

In order to illustrate this point let us analyze the dependence on the anomalous couplings of a subset of the EWPO that contains the $W$ mass ($M_W$), $W$ leptonic width ($\Gamma_{W}^{\ell\nu}$), the $Z$ width into charged leptons ($\Gamma_{\ell\ell}$), the leptonic $Z$ left-right asymmetry ($A_L$), and the invisible $Z$ width ($\Gamma_{inv}^{Z}$). In general we can write the departures of the observables ($\Delta \text{Obs} \equiv \text{Obs}_{\text{SM}} - \text{Obs}_{\text{SM}}$) from the SM predictions as [31]

$$
\begin{pmatrix}
\Delta \Gamma_{\ell\ell} \\
\Delta \Gamma_{inv}^{Z} \\
\Delta A_L \\
\Delta \Gamma_{W}^{\ell\nu}
\end{pmatrix} = M \begin{pmatrix}
f_{1R} \\
f_{1L} \\
f_{3L} \\
f_{\Phi,1}
\end{pmatrix} \frac{v^2}{\Lambda^2}.
$$

(30)

where the matrix $M$ is given by

$$
\begin{pmatrix}
-\frac{4 s^2}{1 - 4 s^2 + 8 s^4} & \frac{2 - 4 s^2}{1 - 4 s^2 + 8 s^4} & \frac{4 s^2 (4 s^2 - 1)}{(c^2 - s^2)(1 - 4 s^2 + 8 s^4)} & -\frac{1 - 2 s^2 - 4 s^4}{2(c^2 - s^2)(1 - 4 s^2 + 8 s^4)} & \frac{4 c s (4 s^2 - 1)}{(c^2 - s^2)(1 - 4 s^2 + 8 s^4)} \\
0 & -2 & 0 & -\frac{1}{2} & 0 \\
\frac{2 s^2 (s^2 - 1/2)^2}{-s^2 + (s^2 - 1/2)^2} & \frac{2 s^2 (s^2 - 1/2)^2}{-s^2 + (s^2 - 1/2)^2} & \frac{-s^2 (s^2 - 1/2)^2}{-s^2 + (s^2 - 1/2)^2} & \frac{-2 s^4 (s^2 - 1/2)^2}{2 (-s^4 + s^2 (s^2 - 1/2)^2)} & \frac{-c s^3}{-s^2 + (s^2 - 1/2)^2} \\
0 & 0 & -\frac{s^2}{c^2 - s^2} & -\frac{c^2}{4(c^2 - s^2)} & -\frac{c s}{c^2 - s^2} \\
0 & 0 & -\frac{3 s^2}{c^2 - s^2} & -\frac{3 c^2}{4(c^2 - s^2)} & -\frac{3 c s}{c^2 - s^2}
\end{pmatrix}.
$$

(31)

It is easy to verify that the matrix $M$ exhibits two zero eigenvalues, indicating that 2 coupling constant combinations can not be determined. In general there are two blind directions even when we consider all LEP observables. In our example, the blind directions are

$$
f_{\Phi,1} = -4 f_{\Phi,1}^{(1)} = -2 f_{\Phi,1}^{(3)} = g^2 f_{BW} \quad \text{and} \quad f_{\Phi,1}^{(3)} = \frac{g^2}{4} f_{BW}.
$$

(32)

\footnote{Here, for the sake of simplicity we assumed lepton flavor universality.}
This means that there are two combinations of operators which do not contribute to these leptonic observables, these are any two linear combination of

\[ O_{\text{lep blind},1} = g^2 (O_{\Phi,1} - \frac{1}{4} \sum_i O_{\Phi, ii}^{(1)} - \frac{1}{2} \sum_i O_{\Phi, ii}^{(2)}) + O_{\text{BW}} \]

\[ O_{\text{lep blind},2} = O_{\text{BW}} + \sum_i O_{\Phi, ii}^{(3)} g^2 \frac{3}{4}. \]

There is a deep relation between operators that do not lead to any tree level contribution to the EWPO and blind directions. In fact, if the elimination of one of these operators using the EOM leads to a combination of operators that contribute at tree level to the EWPO, then this combination defines a blind direction because it has the same S matrix element than the original operator that has no impact on the EWPO [14].

For instance the bosonic operator \( O_{\Phi,2} \) does not contribute to the EWPO since it modifies only the Higgs couplings, therefore, it is a blind operator. Using the EOM given in Eqs. (24)-(26) we can write that

\[ 3g^2 O_{\Phi,2} = 2O_{\text{BW}} + 4O_{W} + 2O_{WW} + \frac{g^2}{2} \sum_i (O_{\Phi, ii}^{(3)} + O_{\Phi, ii})^{(3)} \]

\[ + g^2 \left( \sum_{ij} (y_{ij}^{e}(O_{\Phi, ij}^{e})^{\dagger} + y_{ij}^{u} O_{\Phi, ij}^{u} + y_{ij}^{d}(O_{\Phi, ij}^{d})^{\dagger} + h.c.) \right. \]

The right hand side of the last equation defines a blind direction in the EWPO. In fact, only the operators \( O_{\text{BW}} \) and \( \sum_i O_{\Phi, ii}^{(3)} \) in the right hand side of Eq. (35) contribute to the above leptonic observables, therefore the effect of \( O_{\Phi,2} \) is equivalent to, for these observables,

\[ \frac{2}{3g^2} \left( O_{\text{BW}} + \frac{g^2}{4} \sum_i O_{\Phi, ii}^{(3)} \right) \]

that corresponds to the blind direction in Eq. (35).

The relations (24)-(26) allows the elimination of three dimension–six operators from the basis. As just discussed, in order to take full advantage of the electroweak precision measurements we should avoid the existence of blind directions in the parameter space. This is achieved by using EOM to eliminate operators that contribute at tree level to the EWPO in such way that the new form of the matrix \( M \) in Eq. (30) has a non-vanishing determinant. For example eliminating two combinations of \( O_{\Phi, ii}^{(1)} \) and \( O_{\Phi, ii}^{(3)} \). Furthermore, we also remove \( O_{\Phi,4} \) using Eq. (24). Notice that our choice for the operator basis allows us to take full advantage of the EWPO, as well as, of data on TGV. In brief, the relevant operators for the Higgs physics appearing in our dimension–six operator basis are

\[ \{ O_{GG}, O_{BB}, O_{WW}, O_{BW}, O_{B}, O_{W}, O_{\Phi,2}, O_{\Phi,1}, O_{f\Phi}, O_{\Phi, f}^{(1)}, O_{\Phi, f}^{(3)} \}. \]

except for \( O_{\Phi, L}^{(3)} \). Now we can easily use the available experimental information in order to reduce the number of relevant parameters in the analysis of the Higgs data.

- Taking into account the bulk of precision data on Z and W fermionic currents and oblique corrections discussed above, the coefficients of all operators that modify these couplings are so constrained that they will have no impact in the Higgs physics. Therefore, we will not consider the operators \( (O_{\Phi, f}^{(1)}, O_{\Phi, f}^{(3)}, O_{BW}, O_{\Phi,1}) \) in our analyses.

- Limits on low–energy flavour–changing interactions impose strong bounds on off–diagonal Yukawa couplings [32] [33]. There may still be sizeable flavour changing effects in \( \tau e \) and \( \tau \mu \) [32] [33] which are, however, not relevant to the present analysis. Consequently we also discard from our basis the off–diagonal part of \( O_{\Phi, f} \).

- Flavour diagonal \( O_{f\Phi} \) from first and second generation only affect the present Higgs data via their contribution to the Higgs–gluon–gluon and Higgs–γ–γ vertex at one loop. The loop form factors are very suppressed for light fermions and correspondingly their effect is totally negligible in the analysis. Consequently, we keep only the fermionic operators \( O_{\Phi,33}^{e}, O_{\Phi,33}^{u} \) and \( O_{\Phi,33}^{d} \).
In brief the effective Lagrangian that we use in our analyses is
\[
\mathcal{L}_{\text{eff}} = \frac{\alpha_s v}{8\pi} \frac{f_g^2}{\Lambda^2} O_{GG} + \frac{f_{\Phi,2}}{\Lambda^2} O_{\Phi,2} + \frac{f_{BB}}{\Lambda^2} O_{BB} + \frac{f_{WW}}{\Lambda^2} O_{WW} + \frac{f_{W}}{\Lambda^2} O_{W} + \frac{f_{\tau}}{\Lambda^2} O_{\tau,33} + \frac{f_{\text{bot}}}{\Lambda^2} O_{\text{d},33} + \frac{f_{\text{top}}}{\Lambda^2} O_{\Phi,33} \tag{38}\]

Notice that with this choice of basis all of the dimension-six operators considered contribute to the Higgs–gauge boson and Higgs–fermion couplings at tree level. The tree level information on $htt$ from associate production has very large errors. So quantitatively the effects of the parameter $f_{\text{top}}$ enter via its contribution to the one–loop Higgs couplings to photon pairs and gluon pairs. At present these contributions can be absorbed in the redefinition of the parameters $f_g$ and $f_{WW} + f_{BB}$, therefore, we set $f_{\text{top}} \equiv 0$. In the future, when a larger luminosity will be accumulated, it will be necessary to introduce $f_{\text{top}}$ as one of the parameters in the fit.

### III. ANALYSIS FRAMEWORK

In order to obtain the present constraints on the coefficients of the operators \([38]\) we perform a chi–square test using the available data on the signal strength ($\mu$) from Tevatron, LHC at 7 TeV and 8 TeV for the channels presented in Tables \(\ref{tab:1} \), \(\ref{tab:2} \) and \(\ref{tab:3} \). We will also combine in the chi-square the data coming from the most precise determination of triple electroweak gauge boson couplings, as well as, the one-loop constraints coming from electroweak precision data (EWPD).

In order to predict the expected signal strengths in the presence of the new operators we need to include their effect in the production channels, as well as, in the decay branching ratios. We follow the approach described in \([9]\), which we briefly summarize here. Assuming that the $K$ factor associated with higher order corrections is the same for the SM and new contributions we write

\[
\sigma_Y^{\text{ano}} = \sigma_Y^{\text{ano}} \bigg|_{\text{tree}} \sigma_Y^{\text{SM}} \bigg|_{\text{ano}} \tag{39}\]

\[
\Gamma^{\text{ano}}(h \to X) = \Gamma^{\text{ano}}(h \to X) \bigg|_{\text{tree}} \Gamma^{\text{SM}}(h \to X) \bigg|_{\text{ano}} \tag{40}\]

with the superscripts $\text{ano}$ (SM) standing for the value of the observable considering the anomalous and SM interactions (pure SM contributions). The ratios of the anomalous and SM cross sections of the subprocess $Y$ ($= gg, VBF, VH$ or $ttH$) and of the decay widths are evaluated at tree level, and they are multiplied by the value for the state–of–the–art SM calculations, $\sigma_Y^{\text{SM}}$ and $\Gamma^{\text{SM}}$ presented in Ref. \([53]\). We did not include in our analyses an eventual invisible decay of the Higgs \([54, 55]\), therefore the total width is obtained by summing over the decays into the SM particles. The evaluation of the relevant tree level cross sections was done using the package MadGraph5 \([56]\) with the anomalous Higgs interactions introduced using FeynRules \([57]\). We also cross checked our results using COMPHEP \([58, 59]\) and VBFNLO \([60]\).

For any final state $F$ listed in Tables \(\ref{tab:1} \), \(\ref{tab:2} \) and \(\ref{tab:3} \) we can write the theoretical signal strength as

\[
\mu_F = \frac{\epsilon_{gg}^{\text{ano}}(1 + \xi_g) + \epsilon_{VBF}^{\text{ano}}(1 + \xi_{VBF}) + \epsilon_{VH}^{\text{ano}}(1 + \xi_{VH}) + \epsilon_{\text{ZH}}^{\text{ano}}(1 + \xi_{\text{ZH}}) + \epsilon_{\text{ttH}}^{\text{ano}}}{\epsilon_{gg}^{\text{SM}} + \epsilon_{VBF}^{\text{SM}} + \epsilon_{VH}^{\text{SM}} + \epsilon_{\text{ZH}}^{\text{SM}} + \epsilon_{\text{ttH}}^{\text{SM}}} \times Br^{\text{ano}}[h \to F] \times Br^{\text{SM}}[h \to F] \tag{41}\]

where $\xi_g$, $\xi_{VBF}$ and $\xi_{VH}$ are the pulls associated with the gluon fusion, vector boson fusion and associated production cross section uncertainties (see Eq. \([38]\)), and the branching ratios and the anomalous cross sections are evaluated using the prescriptions \([39]\) and \([40]\). The weight of the different channels to each final state is encoded in the parameters $\epsilon_X$ with $X = VBF, gg, WH, ZH$ and $ttH$.

The search for Higgs decaying into $bb$ pairs takes place through Higgs production in association with a $W$ or a $Z$ so in this case

\[
\epsilon_{gg}^{bb} = \epsilon_{VBF}^{bb} = \epsilon_{VH}^{bb} = \epsilon_{\text{ZH}}^{bb} = \epsilon_{\text{ttH}}^{bb} = 0 , \quad \epsilon_{WW}^{bb} = \epsilon_{ZH}^{bb} = 1 \tag{42}\]

except for the new CMS analysis \([40]\) where

\[
\epsilon_{gg}^{bb} = \epsilon_{VBF}^{bb} = \epsilon_{VH}^{bb} = \epsilon_{\text{ZH}}^{bb} = \epsilon_{\text{ttH}}^{bb} = 0 , \quad \epsilon_{WW}^{bb} = \epsilon_{ZH}^{bb} = 1 \tag{43}\]

is assumed.

The ATLAS and CMS analyses of the $7$ (8) TeV data separate the $\gamma \gamma$ signal into different categories and the contribution of each production mechanism to a given category is presented in Table 6 of ATLAS Ref. \([50]\), Table 1 of ATLAS Ref. \([51]\) and Table 2 of CMS Ref. \([52]\) and we summarized them in Tables \(\ref{tab:4} \) and \(\ref{tab:5} \).
| Channel                  | $\mu^{exp}$          | comment           |
|--------------------------|----------------------|-------------------|
| $pp \to W^+W^-$          | $0.94^{+0.85}_{-0.83}$ | CDF & D0 [38]     |
| $pp \to \tau\tau$       | $1.68^{+2.28}_{-1.68}$ | CDF & D0 [38]     |
| $pp \to bb$              | $1.50^{+0.69}_{-0.72}$ | CDF & D0 [38]     |
| $pp \to \gamma\gamma$   | $5.97^{+3.39}_{-3.12}$ | CDF & D0 [38]     |
| $pp \to \tau\tau$       | $0.7^{+0.7}_{-0.7}$   | ATLAS @ 7 and 8 TeV [39] |
| $pp \to bb$              | $-2.1^{+1.4}_{-1.4}$  | ATLAS @ 7 TeV [40] |
| $pp \to bb$              | $0.6^{+0.7}_{-0.7}$   | ATLAS @ 8 TeV [40] |
| $pp \to ZZ^* \to \ell^+\ell^-\ell^+\ell^-$ | $1.7^{+0.5}_{-0.4}$ | ATLAS @ 7 and 8 TeV [41] |
| $pp \to WW^* \to \ell^+\ell^-\nu\bar{\nu}$ | $0.0^{+0.6}_{-0.6}$ | ATLAS @ 7 TeV [42] |
| $pp \to WW^* \to \ell^+\ell^-\nu\bar{\nu}$ | $1.26^{+0.35}_{-0.35}$ | ATLAS @ 8 TeV [42] |
| $pp \to Z\gamma \to \ell^+\ell^-\gamma$ | $4.7^{+6.89}_{-6.89}$ | ATLAS @ 7 and 8 TeV [43] |
| $pp \to \tau\tau$       | $1.1^{+0.4}_{-0.4}$   | CMS @ 7 and 8 TeV [44] |
| $pp \to bb$              | $1.0^{+0.49}_{-0.49}$ | CMS @ 7 and 8 TeV [45] |
| $pp \to bb$ VBF          | $0.7^{+1.4}_{-1.4}$   | CMS @ 8 TeV [46]  |
| $pp \to ZZ^* \to \ell^+\ell^-\ell^+\ell^-$ | $0.91^{+0.30}_{-0.24}$ | CMS @ 7 and 8 TeV [47] |
| $pp \to WW^* \to \ell^+\ell^-\nu\bar{\nu}$ | $0.91^{+0.44}_{-0.44}$ | CMS @ 7 TeV [48] |
| $pp \to WW^* \to \ell^+\ell^-\nu\bar{\nu}$ | $0.71^{+0.22}_{-0.22}$ | CMS @ 8 TeV [48] |
| $pp \to Z\gamma \to \ell^+\ell^-\gamma$ | $-0.5^{+4.87}_{-4.87}$ | CMS @ 7 and 8 TeV [49] |

Table I: Results included in the analysis for the Higgs decay modes listed except for the $\gamma\gamma$ channels.

| Channel                       | $\mu^{exp}$          | 7 TeV          | 8 TeV          |
|-------------------------------|----------------------|----------------|----------------|
| Unconverted central, low $p_T$ | $0.52^{+1.45}_{-1.40}$ | $0.89^{+0.74}_{-0.71}$ |
| Unconverted central, high $p_T$ | $0.23^{+1.98}_{-1.98}$ | $0.95^{+1.08}_{-0.92}$ |
| Unconverted rest, low $p_T$    | $2.56^{+1.69}_{-1.69}$ | $2.52^{+0.92}_{-0.77}$ |
| Unconverted rest, high $p_T$   | $10.47^{+3.66}_{-3.72}$ | $2.71^{+1.35}_{-1.14}$ |
| Converted central, low $p_T$   | $6.10^{+2.62}_{-2.62}$ | $1.39^{+1.05}_{-0.95}$ |
| Converted central, high $p_T$  | $-4.36^{+1.80}_{-1.80}$ | $2.0^{+1.54}_{-1.26}$ |
| Converted rest, low $p_T$      | $2.73^{+1.98}_{-1.98}$ | $2.22^{+1.57}_{-0.99}$ |
| Converted rest, high $p_T$     | $-1.57^{+2.91}_{-2.91}$ | $1.29^{+1.32}_{-1.26}$ |
| Converted transition          | $0.41^{+3.55}_{-3.66}$ | $2.83^{+1.69}_{-1.69}$ |
| 2-jets / 2-jets high mass tight | $2.73^{+1.92}_{-1.92}$ | $1.63^{+0.83}_{-0.68}$ |
| 2-jets high mass loose        | $-7.77^{+1.99}_{-1.99}$ | $-1.39^{+1.39}$ |
| 2-jets low mass               | $-0.38^{+1.72}_{-1.48}$ | $2.15^{+2.15}$ |
| $E_T^{miss}$ significance     | $2.99^{+2.74}_{-2.15}$ | $2.71^{+2.09}_{-1.66}$ |

Table II: $H \to \gamma\gamma$ results from ATLAS [50, 51] included in our analysis.

With the exception of the above processes, all other channels $F = WW^*, ZZ^*, \bar{\tau}\tau, Z\gamma$ are treated as inclusive,

$$e_{gg}^F = e_{VBF}^F = e_{t\bar{t}H}^F = e_{WH}^F = e_{ZH}^F = 1 .$$ (44)

For some final states the available LHC 8 TeV data has been presented combined with the 7 TeV results. In this case we construct the expected theoretical signal strength as an average of the expected signal strengths for the center–of–mass energies of 7 and 8 TeV. We weight the contributions by the total number of events expected at each energy in the framework of the SM (see Ref. [9] for details).

With all the data described above we perform a $\chi^2$ test assuming that the correlations between the different channels are negligible except for the theoretical uncertainties which are treated with the pull method [62, 63] in order
Table III: $H \rightarrow \gamma\gamma$ results from CMS \cite{52} included in our analysis.

| Channel | $\mu^{\exp}_{\gamma\gamma}$ |
|---------|-----------------|
| $pp \rightarrow \gamma\gamma$ Untagged 3 | 1.48$^{+1.67}_{-1.66}$ $\mu$ |
| $pp \rightarrow \gamma\gamma$ Untagged 2 | 0.024$^{+1.24}_{-1.24}$ $\mu$ |
| $pp \rightarrow \gamma\gamma$ Untagged 1 | 0.194$^{+0.99}_{-0.98}$ $\mu$ |
| $pp \rightarrow \gamma\gamma$ Untagged 0 | 3.83$^{+2.01}_{-1.67}$ $\mu$ |
| $pp \rightarrow \gamma\gamma$ MET | --- $\mu$ |
| $pp \rightarrow \gamma\gamma$ Electron | --- $\mu$ |
| $pp \rightarrow \gamma\gamma$ Muon | --- $\mu$ |

Table IV: Weight of each production mechanism for the different $\gamma\gamma$ categories in the ATLAS analyses of the 7 TeV data (upper values) and 8 TeV (lower values). For the 8 TeV analysis three new exclusive categories enriched in vector boson associated production were added with the 2-jets low mass (lepton tagged) $E_{T}^{\text{miss}}$ significance category being built to select hadronic (leptonic) [invisible] decays of the associated vector boson.

| Channel | $\epsilon_{gg}$ | $\epsilon_{VBF}$ | $\epsilon_{WH}$ | $\epsilon_{ZH}$ | $\epsilon_{t\bar{t}H}$ |
|---------|----------------|----------------|----------------|----------------|----------------|
| Unconverted central, low $p_{T}$ | 1.06 | 0.579 | 0.550 | 0.555 | 0.355 |
| Unconverted central, high $p_{T}$ | 0.760 | 2.27 | 3.03 | 3.16 | 4.26 |
| Unconverted rest, low $p_{T}$ | 1.06 | 0.564 | 0.612 | 0.610 | 0.355 |
| Unconverted rest, high $p_{T}$ | 0.748 | 2.33 | 3.30 | 3.38 | 3.19 |
| Converted central, low $p_{T}$ | 1.06 | 0.578 | 0.581 | 0.555 | 0.357 |
| Converted central, high $p_{T}$ | 0.761 | 2.21 | 3.06 | 3.16 | 4.43 |
| Converted rest, low $p_{T}$ | 1.06 | 0.549 | 0.612 | 0.610 | 0.355 |
| Converted rest, high $p_{T}$ | 0.747 | 2.31 | 3.36 | 3.27 | 3.19 |
| Converted transition | 1.02 | 0.752 | 1.01 | 0.943 | 0.532 |
| 2-jets / 2-jets high mass tight | 0.257 | 11.1 | 0.122 | 0.111 | 0.177 |
| 2-jets high mass loose (only 8 TeV) | 0.514 | 7.74 | 0.160 | 0.170 | 0.171 |
| 2-jets low mass (only 8 TeV) | 0.550 | 0.429 | 9.51 | 9.73 | 3.25 |
| $E_{T}^{miss}$ significance (only 8 TeV) | 0.047 | 0.072 | 11.4 | 26.9 | 20.7 |
| One lepton (only 8 TeV) | 0.025 | 0.086 | 20.2 | 8.71 | 31.9 |

The largest theoretical uncertainties are associated with the gluon fusion subprocess and to account for these errors we introduce two pull factors, one for the Tevatron ($\xi_{g}$) and one for the LHC at 7 and 8 TeV ($\xi^{L}_{g}$). They modify the corresponding predictions as shown in Eq. 41. We consider that the errors associated with the pulls are $\sigma_{g}^{T}=0.43$ and $\sigma_{g}^{L}=0.15$. In the updates including the data released post-Moriond 2013 and the new updates by October 2013, we have also introduced two pull factors to account for theoretical uncertainties.
Table V: Weight of each production mechanism for the different $\gamma\gamma$ categories in the CMS analyses of the 7 TeV data (upper values) and 8 TeV (lower values). $\epsilon_{VH} = \epsilon_{ZH} = \epsilon_{WtH}$. For the $pp \to \gamma\gamma j$ category the 8 TeV data was divided in two independent subsamples labeled as “loose” and “tight” according to the requirement on the minimum transverse momentum of the softer jet and the minimum dijet invariant mass. For the 8 TeV analysis three new exclusive categories were added enriched in vector boson associated production: $\mu$-tag, e-tag and $E_T^{miss}$-tag.

| Channel | $\epsilon_{gg}$ | $\epsilon_{VBF}$ | $\epsilon_{VH}$ | $\epsilon_{WH}$ |
|---------|----------------|-----------------|-----------------|----------------|
| $pp \to \gamma\gamma$ Untagged 3 | 1.04 | 0.637 | 0.808 | 0.355 |
|                                  | 1.06 | 0.558 | 0.675 | 0.343 |
| $pp \to \gamma\gamma$ Untagged 2 | 1.04 | 0.637 | 0.769 | 0.532 |
|                                  | 1.05 | 0.620 | 0.715 | 0.685 |
| $pp \to \gamma\gamma$ Untagged 1 | 1.00 | 0.897 | 1.10 | 0.887 |
|                                  | 0.954 | 1.20 | 1.45 | 1.71 |
| $pp \to \gamma\gamma$ Untagged 0 | 0.702 | 2.43 | 3.69 | 5.50 |
|                                  | 0.833 | 1.66 | 2.66 | 4.45 |
| $pp \to \gamma\gamma jj$ (7 TeV) | 0.036 | 10.5 | 0.118 | 0 |
| $pp \to \gamma\gamma jj$ loose (8 TeV) | 0.535 | 7.31 | 0.348 | 0.856 |
| $pp \to \gamma\gamma jj$ tight (8 TeV) | 0.236 | 11.3 | 0.061 | 0.171 |
| $pp \to \gamma\gamma$, $\mu$-tag (8 TeV) | 0.0 | 0.029 | 16.2 | 35.6 |
| $pp \to \gamma\gamma$, e-tag (8 TeV) | 0.013 | 0.057 | 16.1 | 33.7 |
| $pp \to \gamma\gamma$, $E_T^{miss}$-tag (8 TeV) | 0.241 | 0.358 | 13.2 | 20.2 |

associated with vector boson fusion cross section, one for Tevatron ($\xi_{VBF}^T$) with associated error $\sigma_{VBF}^T = 0.035$ and one for LHC at 7 and 8 TeV ($\xi_{VBF}^L$) with associated error $\sigma_{VBF}^L = 0.03$. Finally theoretical uncertainties from associated production cross section are included with two more pulls, one for Tevatron ($\xi_{VH}^T$) with associated error $\sigma_{VH}^T = 0.075$ and one for LHC at 7 and 8 TeV ($\xi_{VH}^L$) with associated error $\sigma_{VH}^L = 0.05$.

Schematically we can write

$$\chi^2 = \frac{\sum_j (\mu_j - \mu_{j,exp})^2}{\sigma_j^2} + \frac{\sum_{pull} \left( \frac{\xi_{pull}}{\sigma_{pull}} \right)^2}{\rho}$$

(45)

where $j$ stands for channels presented in Tables 11 and 12. We denote the theoretically expected signal as $\mu_{j,exp}$ and error $\sigma_j = \sqrt{\sigma_j^2 + (\sigma_{j,exp})^2}$.

One important approximation in our analyses is that we neglect the effects associated with the distortions of the kinematic distributions of the final states due to the Higgs anomalous couplings arising from their non SM-like Lorentz structure. Thus we implicitly assume that the anomalous contributions have the same detection efficiencies as the SM Higgs. A full simulation of the Higgs anomalous operators taking advantage of their special kinematic features might increase the current sensitivity on the anomalous couplings and it could also allow for breaking degeneracies with those operators which only lead to an overall modification of the strength of the SM vertices (see also [61]), but at present there is not enough public information to perform such analysis outside of the experimental collaborations.

In the next section we will also combine the results of Higgs data from Tevatron and LHC with those from the most precise determination of the triple electroweak gauge boson couplings [27]. For consistency with our multi-parameter analysis, we include the results of the two-dimensional analysis in Ref. [26] which was performed in terms of $\Delta\kappa$, and $\Delta g_1^Z$ with $\Delta\kappa_Z$, $\lambda$, and $\lambda_Z$ as determined by the relations in Eq. (28):

$$\begin{align*}
\kappa &= 0.984^{+0.049}_{-0.049} \\
g_1^Z &= 1.004^{+0.024}_{-0.025}
\end{align*}$$

(46)

with a correlation factor $\rho = 0.11$.

Finally we will account for constraints from electroweak precision data on the higher–order corrections from dimension–six operators in terms of their contribution to the $S, T, U$ parameters as presented for example in Ref. [28]. We will not consider additional effects associated with the possible energy dependence of those corrections. In particular the one–loop contributions from $O_B$, $O_W$, $O_{BB}$, $O_{WW}$ and $O_{s2}$ read
\[
\alpha \Delta S = \frac{1}{6} \frac{e^2}{16 \pi^2} \left\{ 3(f_W + f_B) \frac{m_H^2}{\Lambda^2} \log \left( \frac{\Lambda^2}{m_H^2} \right) + 
+ 2 \left[ (5c^2 - 2)f_W - (5c^2 - 3)f_B \right] \frac{m_Z^2}{\Lambda^2} \log \left( \frac{\Lambda^2}{m_Z^2} \right) 
- \left[ (22c^2 - 1)f_W - (30c^2 + 1)f_B \right] \frac{m_Z^2}{\Lambda^2} \log \left( \frac{\Lambda^2}{m_Z^2} \right) 
- 24(c^2 f_{WW} + s^2 f_{BB}) \frac{m_Z^2}{\Lambda^2} \log \left( \frac{\Lambda^2}{m_H^2} \right) + 2f_{\Phi,2}v^2 \log \left( \frac{\Lambda^2}{m_H^2} \right) \right\},
\]

\[
\alpha \Delta T = \frac{3}{4c^2} \frac{e^2}{16 \pi^2} \left\{ f_B \frac{m_Z^2}{\Lambda^2} \log \left( \frac{\Lambda^2}{m_Z^2} \right) 
+ (c^2 f_W + f_B) \frac{m_Z^2}{\Lambda^2} \log \left( \frac{\Lambda^2}{m_Z^2} \right) 
+ \left[ 2c^2 f_W + (3c^2 - 1)f_B \right] \frac{m_Z^2}{\Lambda^2} \log \left( \frac{\Lambda^2}{m_Z^2} \right) - f_{\Phi,2}v^2 \log \left( \frac{\Lambda^2}{m_H^2} \right) \right\},
\]

\[
\alpha \Delta U = -\frac{1}{3} \frac{e^2s^2}{16 \pi^2} \left\{ -4f_W + 5f_B \frac{m_Z^2}{\Lambda^2} \log \left( \frac{\Lambda^2}{m_Z^2} \right) 
+ (2f_W - 5f_B) \frac{m_Z^2}{\Lambda^2} \log \left( \frac{\Lambda^2}{m_Z^2} \right) \right\}.
\]

(47) (48) (49)

At present the most precise determination of \( S, T, U \) from a global fit yields the following values and correlation matrix

\[
\Delta S = 0.00 \pm 0.10 \quad \Delta T = 0.02 \pm 0.11 \quad \Delta U = 0.03 \pm 0.09
\]

\[
\rho = \begin{pmatrix}
1 & 0.89 & -0.55 \\
0.89 & 1 & -0.8 \\
-0.55 & -0.8 & 1
\end{pmatrix}
\]

(50) (51)

\textbf{IV. PRESENT STATUS}

Initially let us focus on the scenario where we neglect the fermionic Higgs operators, \textit{i.e.} we set \( f_{\text{bot}} = f_\tau = 0 \) and fit the available data using the coefficients of the six bosonic operators \{\( f_g, f_{WW}, f_{BB}, f_{W}, f_B, f_{\Phi,2} \}\} as free independent parameters\(^4\). Considering all Higgs collider (ATLAS, CMS and Tevatron) data, we find \( \chi^2_{\text{min}} = 66.8 \) for the combined analysis and the SM lays at \( \chi^2_M = 58.1 \) \textit{i.e.} within the 3\% CL region in the six dimensional parameter space.

The first column of Figure I displays the chi-square \( \Delta \chi^2 \) dependence upon the six bosonic anomalous couplings after marginalizing over the five undisplayed ones. In this figure we consider all Higgs collider (ATLAS, CMS and Tevatron) data (solid red line), Higgs collider and TGV data (dashed purple line) and Higgs collider, TGV and EWPD (dotted blue line) using \( \Lambda = 10 \text{ TeV} \) in the evaluation of the logarithms in Eqs. \([17, 19]\). As we can see, the \( \Delta \chi^2 \) as a function of \( f_g \) exhibits two degenerate minima in all cases due to the interference between SM and anomalous contributions, that possess exactly the same momentum dependence. Around the secondary minimum the anomalous contribution is approximately twice the SM one but with an opposite sign. The gluon fusion Higgs production cross section is too depleted for \( f_g \) values between the minima, giving rise to the intermediate barrier.

\(^4\) This scenario is a straightforward generalization of the first scenario discussed in Ref. [3].
Figure 1: $\Delta \chi^2$ dependence on the fit parameters when we consider all Higgs collider (ATLAS, CMS and Tevatron) data (solid red line), Higgs collider and TGV data (dashed purple line) and Higgs collider, TGV and EWP data (dotted blue line). The rows depict the $\Delta \chi^2$ dependence with respect to the fit parameter shown on the left of the row with the anomalous couplings $f/\Lambda^2$ given in TeV$^{-2}$. In the first column we use $f_g, f_{WW}, f_{BB}, f_W, f_B,$ and $f_{\Phi,2}$ as fit parameters with $f_{bot} = f_\tau = 0$. In the second column the fitting parameters are $f_g, f_{WW} = -f_{BB}, f_W, f_B, f_{\Phi,2},$ and $f_{bot}$ with $f_\tau = 0$. In the panels of the right column we fit the data in terms of $f_g, f_{WW} = -f_{BB}, f_W, f_B, f_{\Phi,2}, f_{bot},$ and $f_\tau$. 
Figure 2: Chi–square dependence on the Higgs branching ratios (left panels) and production cross sections (right panels) when we consider all Higgs collider and TGV data. In the upper panels we used \( f_g, f_{WW}, f_{BB}, f_{W}, f_{B}, \) and \( f_{\Phi,2} \) as fitting parameters with \( f_{\text{bot}} = f_{\tau} = 0 \), while in the middle panels the fitting parameters are \( f_g, f_{WW} = -f_{BB}, f_{W}, f_{B}, f_{\Phi,2}, \) and \( f_{\text{bot}} \) with \( f_{\tau} = 0 \). In the lower row we parametrize the data in terms of \( f_g, f_{WW} = -f_{BB}, f_{W}, f_{B}, f_{\Phi,2}, f_{\text{bot}}, \) and \( f_{\tau} \). The dependence of \( \Delta \chi^2 \) on the branching ratio to the fermions not considered in the analysis arises from the effect of the other parameters in the total decay width.

\( f_B \) and \( f_W \) are the only fit parameters that modify the triple gauge vertices at tree level, therefore, they are the ones that show the largest impact of the TGV data as can be seen in the corresponding panels of Fig. 1. \( f_W \) is the most constrained parameter by the inclusion of the TGV data since this is the only one that modifies the most precisely determined TGV \( g_1^2 \). Moreover, the inclusion of the EWPD in the fit reduces significantly the errors on \( f_B, f_{\Phi,2}, f_{BB}, \) and \( f_{WW} \). As expected, there is little impact on \( f_g \) from the addition of the TGV and EWPD data sets.

The best fit values and 90% CL allowed ranges for the couplings and observables in the combined analysis of Higgs collider and TGV data can be found in the two first columns in Table VI. Inclusion of the TGV data has almost no quantitative impact on the values of \( \chi^2_{\text{min}} \) nor the SM CL. Adding the EWPD increases \( \chi^2_{\text{min,SM}} \) to 67.9 (69.9) so the SM lies in the full combined analysis at the 9% CL six–dimensional region in agreement with these combined results at 0.1\( \sigma \) level.

We translate the results displayed in Fig. 1 in terms of physical observables in Fig. 2 which shows the \( \Delta \chi^2 \) dependence.
on the Higgs decay branching ratios and production cross sections\(^5\). As we can see from the two top panels, the SM predictions are within the 68% CL allowed ranges using the Higgs collider data together with TGV data. Notice that with the presently available data the Higgs branching ratios are known with a precision around 20% while the production cross sections are known with an uncertainty of 30%.

We depict in Figure 3 the 95% and 99% CL allowed regions of the plane \(f_{WW} \times f_{BB}\), after marginalizing over the undisplayed variables, when we consider only the Higgs collider data. As we can see there is a strong anti-correlation between \(f_{WW}\) and \(f_{BB}\) since they are the dominant contributions to the Higgs branching ratio into two photons which is proportional to \(f_{WW} + f_{BB}\). The 95% CL allowed region is formed by two narrow islands: one with small departures from the SM contribution and a second one around the anomalous couplings values such that their contribution to the Higgs decay amplitude to photons is twice the SM one but with the opposite sign. This degeneracy of the minima is not exact since the \(f_{WW}\) and \(f_{BB}\) couplings not only contribute to Higgs decay into photons, but also to its decay into \(WW^*\) and \(ZZ^*\) as well as in \(Vh\) associated and vector boson fusion production mechanisms, lifting the degeneracy of the local minima. Notice also that after marginalization over \(f_{BB}\) \((f_{WW})\), the one-dimensional \(\Delta \chi^2\) curve for \(f_{WW}\) \((f_{BB})\) shown in second (third) row of the first column in Fig. 1 has only one minima and the anti-correlation is translated in these two curves being close to mirror symmetric.

![Figure 3](image-url)

**Figure 3:** We display the 95% and 99% CL allowed regions in the plane \(f_{WW} \times f_{BB}\) when we fit the Higgs collider data varying \(f_g, f_{WW}, f_{BB}, f_w, f_B,\) and \(f_{\Phi,2}\). The star stands for the global minima and we marginalized over the undisplayed parameters.

Figure 4 contains the 68%, 90%, 95%, and 99% CL 2-dimensional projection in the plane \(f_g \times f_{\Phi,2}\) after marginalization over the four undisplayed parameters. The results are shown for the combination of Higgs collider and TGV data sets. As we can see, this figure exhibits two isolated islands that originate from the interference between anomalous and SM contributions to the Higgs coupling to two gluons. Within each island there is an anti-correlation between \(f_g\) and \(f_{\Phi,2}\) that stems from the fact that the anomalous contribution to the Higgs gluon fusion production is proportional to \(F_{99}^{SM} f_{\Phi,2} + 2 f_g\) where \(F_{99}^{SM} \simeq 0.7\) is the SM loop contribution to the \(Hg\) vertex.

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\(^5\) Here we do not include EWPD to avoid the introduction of a model dependent scale needed to evaluate the logarithms present in Eqs. (47)–(49).
Figure 4: We present the 68%, 90%, 95%, and 99% CL allowed regions in the plane $f_g \times f_{\Phi,2}$ when we fit the Higgs collider and TGV data varying $f_g$, $f_{WW}$, $f_{BB}$, $f_W$, $f_B$, and $f_{\Phi,2}$. The stars stand for the global minima and we marginalized over the undisplayed parameters.

The left panel of Fig. 5 displays the correlations between the Higgs branching ratio into photons and its gluon fusion production cross section in the scenario where with $f_{\text{bot}} = f_r = 0$. Clearly, these quantities are anti-correlated since their product is the major source of Higgs events decaying into two photons.

Let us now turn to the effects of including the fermionic operators in the analysis. We first do so by augmenting the set of parameters by the anomalous bottom Yukawa-like coupling $f_{\text{bot}}$, however, to simplify the numerical analyses we enforce the strong correlation observed in Fig. 3 between $f_{WW}$ and $f_{BB}$ imposing that $f_{WW} = - f_{BB}$. Therefore, our free parameters are $\{f_g, f_W, f_B, f_{WW} = - f_{BB}, f_{\Phi,2}, f_{\text{bot}}\}$, where we are still keeping $f_r = 0$.

We present in the middle panels of Fig. 1 the chi-square as a function of the fitting parameters in this case. First we see that the $\Delta \chi^2$ dependence of $f_{\text{bot}}$ presents two degenerate minima, one small correction to the SM Yukawa coupling and one larger positive which will flip the sign of the $Hbb$ coupling. Comparing with the first column of panels in this figure we see that the allowed range for $f_g$ becomes much larger and the one for $f_{\text{bot}}$ is also large. This behavior emanates from the fact that at large $f_{\text{bot}}$ the Higgs branching ratio into $b$–quark pairs approaches 1, so to fit the data for any channel $F \neq bb$, the gluon fusion cross section must be enhanced in order to compensate the dilution of the $H \rightarrow F$ branching ratios. This is clearly shown in Fig. 6 which depicts the strong correlation between the allowed values of $f_{\text{bot}} \times f_g$. This correlation has an impact on the determination of the gluon fusion production cross section and the Higgs branching ratio into photon pairs as illustrated in the right panel of Fig. 5 which shows that the gluon fusion production cross section can now be much larger than the SM cross section but only as long as the Higgs branching ratio into photons is below the SM value in order to fit the observed rate of $\gamma \gamma$ events. On the other hand allowing for $f_{\text{bot}} \neq 0$ has a small impact on the parameters affecting the Higgs couplings to electroweak gauge bosons $f_W$, $f_B$, $f_{WW} = - f_{BB}$, and $f_{\Phi,2}$ as seen by comparing the corresponding left and central panels of Fig. 1 even prior to the inclusion of TGV constraints on $f_W$ and $f_B$.

The effect of $f_{\text{bot}}$ can also be understood by comparing the upper and central lines in Fig. 2 which contain the chi-square dependence on Higgs branching ratios (left) and production cross sections (right) for the analysis with $f_{\text{bot}} = 0$ (upper) and $f_{\text{bot}} \neq 0$ (central). We can immediately see that the bounds on branching ratios and cross sections get loosened, with the VBF and VH production cross sections being the least affected quantities while the
Figure 5: In the left (right) panel we present the 68%, 90%, 95%, and 99% CL allowed regions in the plane $\sigma_{\text{ano}}^{gg}/\sigma_{\text{SM}}^{gg} \times \text{Br}(h \rightarrow \gamma\gamma)_{\text{ano}}/\text{Br}(h \rightarrow \gamma\gamma)_{\text{SM}}$ when we fit the Higgs collider and TGV data varying $f_g$, $f_{WW}$, $f_{BB}$, $f_W$, $f_B$, and $f_{\Phi,2}$ ($f_g$, $f_{WW} = -f_{BB}$, $f_W$, $f_B$, $f_{\Phi,2}$, and $f_{\text{bot}}$). The stars stand for the global minima and we marginalized over the undisplayed parameters.

The gluon fusion cross section is the one becoming less constrained. The reason for this deterioration of the constraints is due to the strong correlation between $f_g$ and $f_{\text{bot}}$ we just mentioned.

The impact of $f_{\text{bot}}$ on the fit is due to the absence of data on the direct process $pp \rightarrow h \rightarrow b\bar{b}$ due to the huge SM backgrounds. One way to mitigate the lack of information on this channel is to have smaller statistical errors in the processes taking place via VBF or VH associated production. However, this will require a larger data sample than that which is presently available.

Finally we study the effect of allowing $f_\tau \neq 0$. For the sake of simplicity we keep the number of free parameters equal to seven and we choose them to be $\{f_g, f_W, f_B, f_{WW}, f_{BB}, f_{\Phi,2}, f_{\text{bot}}, f_\tau\}$. The data, once more, the strong correlation between $f_{WW}$ and $f_{BB}$ to reduce the number of free parameters to a treatable level. We present in the right panels of Fig. 1 the chi–square as a function of the free parameters in this case and in the lower panels of Fig. 2 the corresponding dependence for the decay branching ratios and production cross sections. The results are that the inclusion of $f_\tau$ in the analysis does not introduce any further strong correlation. This is so because the data on $pp \rightarrow h \rightarrow \tau^+\tau^-$ cuts off any strong correlation between $f_\tau$ and $f_g$. The determination of the parameters is not affected very much with respect to the previous case with $f_{\text{bot}} \neq 0$ and $f_\tau = 0$. Concerning the observables only $\tau\tau$ Higgs branching ratio is affected. The corresponding best fit values and allowed 90% CL ranges for the parameters and observables are given in the right two columns in Table VI.

We see that at the best fit point the present global analysis favors a $BR^{\text{ano}}_{\tau\tau}/BR^{\text{SM}}_{\tau\tau}$ bigger than 1 (1.1) which leads to two possible values of $f_\tau$ one small positive correction to the negative SM Yukawa coupling and one larger positive which will flip the sign of the $H\tau\tau$ coupling but give the same absolute value. This is the origin of the two minima observed in the lowest panel in Fig. 1. Also, the inclusion of the fermion couplings has no impact on the values of $\chi^2_{\text{min},(SM)}$, and it still holds that the SM is in overall agreement with the Higgs and TGV results at better than 9%
Figure 6: We present the 68%, 90%, 95%, and 99% CL allowed regions in the plane \( f_{\text{bot}} \times f_g \) when we fit the Higgs collider and TGV data varying \( f_g, f_W, f_B, f_{WW} = -f_{BB}, f_{k,2}, \) and \( f_{\text{bot}} \). The stars stand for the global minima and we have marginalized over the undisplayed parameters.

V. DISCUSSIONS AND CONCLUSIONS

As the ATLAS and CMS experiments accumulate more and more luminosity we start to better probe the couplings of the recently discovered “Higgs–like” state. In this work we used a bottom–up approach to describe departures of the Higgs couplings from the SM predictions. In a model independent framework these effects can be parametrized in terms of an effective Lagrangian. Assuming that the observed state is a member of an \( SU(2)_L \) doublet, and therefore the \( SU(2)_L \times U(1)_Y \) gauge symmetry is linearly realized, they appear at lowest order as dimension–six operators with unknown coefficients containing all the SM fields including the light scalar doublet; for details return to Sec. II where we give the full list of operators affecting the Higgs couplings to gauge bosons and fermions. Not all the operators in Eqs. (2) and (17) are independent because at any order they are related by the equations of motion. This allows for a “freedom of choice” in the election of the basis of operators to be used in the analysis.

We have argued in Sec. II C that in the absence of any \textit{a priori} knowledge on the form of the new physics the most sensible choice of basis should contain operators whose coefficients are more easily related to existing data from other well tested sectors of the theory, \textit{i.e.}, not only the LHC data on the Higgs production, but also EWPD and searches for anomalous triple gauge vertices. In this approach we reduce the operator basis to the nine operators in Eq. (38) whose coefficients are still not severely constrained by non–Higgs observables, and which are directly testable with an analysis of the existing Higgs data. The summary of our present determination of Higgs couplings, production cross sections and decay branching ratios from the analysis of the Higgs and TGV data can be found in Table VI.

Generically in any analysis, we obtained that the SM predictions for each individual coupling and observable are well within the corresponding 68% ranges.

The presence of non–vanishing coefficients for the dimension–six operators alters the high energy behavior of the scattering amplitudes of SM particles. The scale where unitarity is violated at tree level in a given process can
be used as a rough estimation for the onset of new physics. For instance, the $2 \to 2$ scattering of SM (Higgs or gauge) bosons has been used to test the validity of a theory containing dimension–six effective operators [64, 65]. The operators $O_W$ and $O_B$ give rise to a contribution to the neutral $W_L^+ W_L^-$, $Z_L Z_L$ and $HH$ channels which grows like $(f_{W,B} s / M_W^2)^2$ [64]. Taking the largest value of the 90\% CL regions a study of unitarity violation indicates that the scale of New Physics beyond SM is larger than $\simeq 2$ TeV; a result in agreement with the EWPD.

Acknowledgments

We thank E. Masso for discussions and participation on the early stages of this work. O.J.P.E is grateful to the Institute de Physique Théorique de Saclay for its hospitality. O.J.P.E. is supported in part by Conselho Nacional de Desenvolvimento Científico e Tecnológico (CNPq) and by Fundação de Amparo à Pesquisa do Estado de São Paulo (FAPESP); M.C.G-G is also supported by USA-NSF grant PHY-09-6739, by CUR Generalitat de Catalunya grant 2009SGR502 and together with J.G-F by MICINN FPA2010-20807 and consolidador-ingenio 2010 program CSD-2008-0037 and by EU grant FP7 ITN INVISIBLES (Marie Curie Actions PITN-GA-2011-289442). J.G-F is further supported by Spanish ME FPU grant AP2009-2546. T.C is supported by USA-NSF grant PHY-09-6739.

\[ \begin{array}{|c|c|c|c|}
\hline
& \text{Fit with } f_{\text{tot}} = f_r = 0 & \text{Fit with } f_{\text{tot}} \text{ and } f_r & \\
\hline
f_{\phi}/\Lambda^2 (\text{TeV}^{-2}) & 1.2, 22 & [-3.3, 5.1] \cup [19, 26] & 2.1, 21 & [-5.3, 5.8] \cup [17, 22] \\
\hline
f_{WW}/\Lambda^2 (\text{TeV}^{-2}) & 1.5 & [-3.2, 8.2] & 0.65 & [-4.2, 7.7] \\
\hline
f_{BB}/\Lambda^2 (\text{TeV}^{-2}) & -1.6 & [-7.5, 5.3] & -0.65 & [-7.7, 4.2] \\
\hline
f_{WW}/\Lambda^2 (\text{TeV}^{-2}) & 2.1 & [-5.6, 9.6] & 1.7 & [-5.4, 9.8] \\
\hline
f_{BB}/\Lambda^2 (\text{TeV}^{-2}) & -10 & [-29, 8.9] & -7.9 & [-28, 11] \\
\hline
f_{BB}/\Lambda^2 (\text{TeV}^{-2}) & -1.0 & [-10, 8.5] & -1.3 & [-9.8, 7.5] \\
\hline
f_{\text{tot}}/\Lambda^2 (\text{TeV}^{-2}) & --- & --- & 0.01, 0.84 & [-0.28, 0.24] \cup [0.55, 1.3] \\
\hline
f_r/\Lambda^2 (\text{TeV}^{-2}) & --- & --- & -0.01, 0.37 & [-0.07, 0.05] \cup [0.26, 0.49] \\
\hline
BR_{WW}^{s} / BR_{WW}^{SM} & 1.2 & [0.78, 1.7] & 1.2 & [0.55, 1.9] \\
\hline
BR_{WW}^{t} / BR_{WW}^{SM} & 1.0 & [0.89, 1.1] & 1.2 & [0.51, 1.9] \\
\hline
BR_{WW}^{b} / BR_{WW}^{SM} & 1.2 & [0.84, 1.5] & 1.4 & [0.62, 2.2] \\
\hline
BR_{WW}^{t} / BR_{WW}^{SM} & 1.0 & [0.92, 1.1] & 0.89 & [0.46, 1.3] \\
\hline
\sigma_{bb}^{W,SM} / \sigma_{bb}^{W} & 1.0 & [0.92, 1.1] & 1.1 & [0.42, 2.6] \\
\hline
\sigma_{bb}^{W,SM} / \sigma_{bb}^{W} & 0.88 & [0.59, 1.3] & 0.73 & [0.38, 2.0] \\
\hline
\sigma_{bb}^{W,SM} / \sigma_{bb}^{W} & 1.1 & [0.52, 1.9] & 1.1 & [0.58, 1.8] \\
\hline
\sigma_{bb}^{W,SM} / \sigma_{bb}^{W} & 0.82 & [0.43, 1.4] & 0.96 & [0.47, 1.5] \\
\hline
\end{array} \]

Table VI: Best fit values and 90\% CL allowed ranges for the combination of all available Tevatron and LHC Higgs data as well as TGV.
M. Perez-Victoria, Phys. Rev. D 78, 013010 (2008) [arXiv:0803.4008 [hep-ph]].
[32] S. Kanemura, T. Ota and K. Tsumura, Phys. Rev. D 73 (2006) 016006 [hep-ph/0505191].
[33] P. Paradisi, JHEP 0602 (2006) 050 [hep-ph/0508054].
[34] E. Gabrielli and B. Mele, Phys. Rev. D 83 (2011) 073009 [arXiv:1102.3361 [hep-ph]].
[35] S. Davidson and G. J. Grenier, Phys. Rev. D 81 (2010) 095016 [arXiv:1001.0434]; S. Davidson and P. Verdirer, [arXiv:1211.1248 [hep-ph]].
[36] A. Goudelis, O. Lebedev and J. -h. Park, Phys. Lett. B 707 (2012) 369 [arXiv:1111.1715].
[37] G. Blankenburg, J. Ellis and G. Isidori, Phys. Lett. B 712, 386 (2012) [arXiv:1202.5704 [hep-ph]].
[38] B. Tuchming [D0 and CDF Collaborations], [arXiv:1307.4873 [hep-ex]].
[39] ATLAS Collaboration, ATLAS-CONF-2012-160.
[40] ATLAS Collaboration, ATLAS-CONF-2013-079.
[41] ATLAS Collaboration, ATLAS-CONF-2013-013.
[42] ATLAS Collaboration, ATLAS-CONF-2013-030.
[43] ATLAS Collaboration, ATLAS-CONF-2013-009.
[44] CMS Collaboration, CMS PAS HIG-13-004.
[45] CMS Collaboration, CMS PAS HIG-13-012.
[46] CMS Collaboration, CMS PAS HIG-13-011.
[47] CMS Collaboration, CMS PAS HIG-13-002.
[48] CMS Collaboration, CMS PAS HIG-13-003.
[49] S. Chatrchyan et al. [CMS Collaboration], [arXiv:1307.5515 [hep-ex]].
[50] ATLAS Collaboration, ATLAS-CONF-2012-091.
[51] ATLAS Collaboration, ATLAS-CONF-2013-012.
[52] CMS Collaboration, CMS PAS HIG-13-001.
[53] LHC Higgs Cross Section Working Group, S. Dittmaier et al., (2011), [arXiv:1101.0593].
[54] J. R. Espinosa, M. Muhlleitner, C. Grojean, and M. Trott, (2012), [arXiv:1205.6790].
[55] M. Raidal and A. Strumia, Phys.Rev. D84, 077701 (2011), [arXiv:1108.4903].
[56] J. Alwall, M. Herquet, F. Maltoni, O. Mattelaer, and T. Stelzer, JHEP 1106, 128 (2011), [arXiv:1106.0522].
[57] N. D. Christensen and C. Duhr, Comput.Phys.Commun. 180, 1614 (2009), [arXiv:0806.4194].
[58] A. Pukhov et al., (1999), [arXiv:hep-ph/9908288].
[59] CompHEP Collaboration, E. Boos et al., Nucl.Instrum.Meth. A534, 250 (2004), [arXiv:hep-ph/0403113].
[60] K. Arnold et al., (2011), [arXiv:1107.4038].
[61] E. Masso and V. Sanz, [arXiv:1211.1320 [hep-ph]].
[62] G. Fogli, E. Lisi, A. Marrone, D. Montanino, and A. Palazzo, Phys.Rev. D66, 053010 (2002), [arXiv:hep-ph/0206162].
[63] M. Gonzalez-Garcia and M. Maltoni, Phys.Rept. 460, 1 (2008), [arXiv:0704.1800].
Phys. Rev. D86 (2012) 035015 [arXiv:1206.0535 [hep-ph]]; JHEP 1207 (2012) 175 [arXiv:1205.8542 [hep-ph]].
[64] G. J. Gounaris, J. Layssac, J. E. Paschalidis and F. M. Renard, Z. Phys. C 66, 619 (1995) [hep-ph/9409260].
[65] G. J. Gounaris, J. Layssac and F. M. Renard, Phys. Lett. B 332, 146 (1994) [hep-ph/9311370].