Chapter 2
Chiral EW Lagrangian

The ATLAS and CMS collaborations at LHC have found a new boson compatible with the SM Higgs [1, 2], with a mass of nearly.[1] $M_h \simeq 125\,\text{GeV}$. Furthermore, the most probable J^P quantum numbers are 0^+, and couplings with other particles are in agreement with the SM Higgs, although with moderate precision. Moreover, there is a mass gap for the presence of new physics [5–7] until an energy of about 600–700 GeV, or even higher for the presence of new vector resonances.

However, the data is still compatible with either an elementary or a composite Higgs: this last possibility will be considered in this work. The mass gap between the M_W, M_Z and M_h masses, all of $\mathcal{O}(100\,\text{GeV})$, and the new physics scale (if there is one within reach), suggests that the Higgs boson and the would-be Goldstone bosons $\omega^\pm$ and z could be (pseudo) Goldstone Boson [8–15], related with a global spontaneous symmetry breaking extending the $SU(2)_L \times SU(2)_R \to SU(2)_{L+R}$ global symmetry breaking of the SM. There are several models with specific implementations for the relevant global symmetry breaking pattern: the (Minimal) Composite Higgs Model based on the coset SO(5)/SO(4) [16–19], dilaton models [20, 21] and others [22].

The old electroweak chiral Lagrangian (ECL) [23–40], based on standard chiral perturbation theory (ChPT) of QCD [41–44], assumed a Higgsless model, but solved the problems that the Higgs was intended for by supposing a strongly interacting regime for the EWSBS instead. In view that the h has been found, it can be extended to include the new Higgs-like particle found at the LHC (refs. [45–56]). One of the goals of this work will be to expose this extension, considering non-linear Electroweak Chiral Lagrangians as a low-energy (M_W, $M_h \ll \sqrt{s} \ll 3\,\text{TeV}$) parameterization of the new physics at the TeV scale.

[1] $M_h^{\text{ATLAS}} = 125.5 \pm 0.6\,\text{GeV}$ and $M_h^{\text{CMS}} = 125.7 \pm 0.4\,\text{GeV}$, according to [3] and [4], respectively.

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2.1 Equivalence Theorem

To simplify the computations, we will make use of the Equivalence Theorem (ET) [57–60], which states that, in the regime $s \gg M_h^2, M_W^2, M_Z^2 \simeq (100\,\text{GeV})^2$, we can identify the longitudinal modes of gauge bosons with the would-be Goldstones[2] (R_ξ gauge). For example,

$$T(W_L^a W_L^b \to W_L^c W_L^d) = T(\omega^a \omega^b \to \omega^c \omega^d) + O\left(\frac{M_W}{\sqrt{s}}\right), \tag{2.1}$$

where T stands for the corresponding scattering amplitude. In fact, this theorem can be interpreted as if, at sufficiently large energies, the symmetry $SU(2)_L \times U(1)_Y$ was not spontaneously broken, so that the three would-be Goldstones coming from the broken $SU(2)_L \times U(1)_Y$ were directly observable as scalar physical particles, which indeed would correspond to the longitudinal modes of the gauge bosons. So, the non-gauged (but with the broken symmetry $SU(2)_L \times U(1)_Y$) Lagrangian can be used directly to compute the scattering amplitudes. At lower energies, these would-be Goldstones give rise to the longitudinal modes of gauge bosons through a rotation in the coordinates (gauge), according to the Higgs mechanism explained, for instance, in Ref. [61] and Sect. 1.

Let us illustrate the application of the ET. We will compare the (exact) tree level computation and the Equivalence Theorem for WW and ZZ scattering on the SM, which can be found in Refs. [62, 63]. The complete SM tree level matrix element for $W_L^+ W_L^- \to HH$ is

$$\begin{aligned}\mathcal{A} = \frac{g^2}{4-(1-4M_W^2/s)} &\left\{ \frac{M_h^2}{M_W^2}\left[1 + \frac{3M_h^2}{s-M_h^2} + M_h^2\left(\frac{1}{t-M_W^2} + \frac{1}{u-M_W^2}\right)\right]\right.\\ &+2\left[1 - \frac{9M_h^2}{s-M_h^2} + 4\frac{M_W^2}{s}\left(1+\frac{3M_h^2}{s-M_h^2}\right)\right]\\ &\left.+2\left[s - 2M_h^2 - 4M_W^2 + 8\frac{M_W^4}{s}\right]\left(\frac{1}{t-M_W^2} + \frac{1}{u-M_W^2}\right)\right\},\end{aligned} \tag{2.2}$$

whereas, for the ET $\omega^+\omega^- \to HH$,

$$\begin{aligned}\tilde{\mathcal{A}} = \frac{g^2}{4}&\left\{ \frac{M_h^2}{M_W^2}\left[1 + \frac{3M_h^2}{s-M_h^2} + M_h^2\left(\frac{1}{t-M_W^2} + \frac{1}{u-M_W^2}\right)\right]\right.\\ &\left.+2\left[1 + (s-M_h^2)\left(\frac{1}{t-M_W^2} + \frac{1}{u-M_W^2}\right)\right]\right\}.\end{aligned} \tag{2.3}$$

[2] See app. D.5 for a brief historical review about the discussion concerning the hypothesis of the ET.

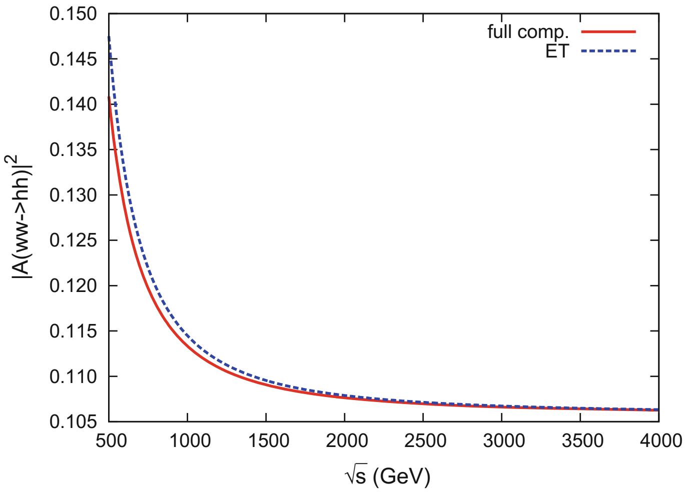

Fig. 2.1 Comparison between the full LO scattering amplitude for $\omega\omega \to hh$ (Eq. 2.2 and *green dashed line* on the plot) and that computed though the equivalence theorem (Eq. 2.3 and *red solid line* on the plot). See Ref. [62, 63] for the computation. $x = \cos\theta = 0.3$

Both Eqs. 2.2 and 2.3, can be evaluated for certain values of the scattering angle θ and s. According to the Equivalence Theorem, both results should converge in the limit $(M_W^2/s) \to 0$. In order to recover $u = u(s, \theta)$ and $t = t(s, \theta)$ as a function of the scattering angle θ and the squared center of mass energy s, see the expressions of appendix A.

In collaboration with Prof. Stefano Moretti (University of Southampton), we have tested the equivalence theorem in this way (see Fig. 2.1). The aim of this test was to cross-check a modified version of the Monte Carlo (MC) program MadGraph [64]. According to this experience, when dealing with expressions from other authors or Monte Carlo (MC) programs, it is crucial that both the masses and the coupling are compatible. Sometimes, the Monte Carlo program accepts masses and couplings separately, without ensuring compatibility. Or you can develop a program for generating points in the phase space that, of course, will take the masses of gauge bosons as an input. It can also happen that the MC program uses the so-called *Complex Mass Scheme*, which would also require to modify the couplings. The fact is that the cancellation between diagrams which leads to a weakly interacting EWSBS on the SM can be very easily spoiled at TeV energies because of using incompatible values for the constants. Thus, the set of numerical values of the LO couplings should verify, with high precision, the well-known SM relations.[3]

[3] See, for instance, Ref. [65] or [61].

2.2 The Chiral Lagrangian and Its Parameterizations

In Ref. [66] we present the effective Lagrangian describing the low-energy dynamics of four light modes: three would-be Goldstone Bosons ω^a (WBGBs) and the Higgs-like particle[4] h. This model is valid for the energy range $M_h, M_W, M_Z \simeq (100\,\text{GeV})^2 \ll s \ll 4\pi v \simeq 3\,\text{TeV}$. The effective Lagrangian is

$$\mathcal{L} = \frac{v^2}{4} g(h/f) \text{Tr}[(D_\mu U)^\dagger D^\mu U] + \frac{1}{2}\partial_\mu h \partial^\mu h - V(h), \tag{2.4}$$

where $v = 246\,\text{GeV}$ is the SM Higgs-doublet vacuum expectation value; f, a new dynamical energy scale; and $g(x)$, an arbitrary analytical funtion of the scalar field,

$$g(h/f) = 1 + \sum_{n=1}^{\infty} g_n \left(\frac{h}{f}\right)^n = 1 + 2\alpha\frac{h}{f} + \beta\left(\frac{h}{f}\right)^2 + \cdots \tag{2.5}$$

Instead of using a power expansion over $1/f$, we could also choose $1/v$, v being the vacuum expectation value,

$$g(h/v) = 1 + 2a\frac{h}{v} + b\left(\frac{h}{v}\right)^2 + \cdots . \tag{2.6}$$

V is an arbitrary analytical potential for the scalar field,

$$V(h) = \sum_{n=0}^{\infty} V_n h^n \equiv V_0 + \frac{M_h^2}{2} h^2 + \sum_{n=3}^{\infty} \lambda_n h^n \tag{2.7}$$

U is a field taking values in the $SU(2)$ coset. In this work, unless otherwise stated, we will use the so-called *spherical parameterization*,

$$U = \sqrt{1 - \frac{\tilde{\omega}^2}{v^2}} + i\frac{\tilde{\omega}}{v}, \tag{2.8}$$

$\tilde{\omega} = \omega_a \tau^a$ being the would-be Goldstone bosons (WBGB) field and τ^a, the Pauli matrices. Note the presence of the non-linear term $\sqrt{1 - (\tilde{\omega}^2/v^2)}$. This is the main difference from linear approaches like [45–56].

The covariant derivative of the U field (Eq. 2.8) is defined as

$$D_\mu U = \partial_\mu U + i\hat{W}_\mu U - iU\hat{B}_\mu, \tag{2.9}$$

where

[4] Also called φ in some early works like [66].

$$\hat{W}_\mu = g W_{\mu,i} \frac{\tau^i}{2}, \qquad \hat{B}_\mu = g' B_\mu \frac{\tau^3}{2} \tag{2.10a}$$

$$\hat{W}_{\mu\nu} = \partial_\mu \hat{W}_\nu - \partial_\nu \hat{W}_\mu + i[\hat{W}_\mu, \hat{W}_\nu], \qquad \hat{B}_{\mu\nu} = \partial_\mu \hat{B}_\nu - \partial_\nu \hat{B}_\mu. \tag{2.10b}$$

We follow the chiral counting of refs. [38, 41, 53]. Alonso et al. [51] have also studied the counting of Electroweak Chiral Lagrangians, but from a different point of view. Note that the chiral counting which we use is explained in detail in our Ref. [67], and is applied only once the approximation $M_h^2, M_W^2, M_Z^2 \ll s$ is taken into account.

Anyway, the 'chiral counting' involves organizing the invariant terms of the Effective Lagrangian by means of their *chiral dimension*. That is, a term $\mathcal{L}_d$ with *chiral dimension* d will contribute to $\mathcal{O}(p^d)$ in the corresponding power momentum expansion. The derivatives and the masses of the dynamical particles (when they are not neglected) are considered as soft scales of the Effective Theory, of order $\mathcal{O}(p)$. To sum up,

$$\partial_\mu,\ M_W,\ M_Z,\ M_h \sim \mathcal{O}(p) \tag{2.11a}$$

$$D_\mu U,\ V_\mu,\ g' v \mathcal{T},\ \hat{W}_\mu,\ \hat{B}_\mu \sim \mathcal{O}(p) \tag{2.11b}$$

$$\hat{W}_{\mu\nu}\ ,\hat{B}_{\mu\nu} \sim \mathcal{O}(p^2) \tag{2.11c}$$

The parameters a and b can be adjusted to fit different theoretical models, and the NLO parameters will depend on the renormalization of the underlying theory. For instance,

- $a^2 = b = 0$ .... Higgsless ECL (ruled out) [23, 25, 26],
- $a^2 = b = 1$ .... SM,
- $a^2 = 1 - \frac{v^2}{f^2}$, $b = 1 - \frac{2v^2}{f^2}$ .... $SO(5)/SO(4)$ MCHM [16–19],
- $a^2 = b = \frac{v^2}{\hat{f}^2}$ .... Dilaton [20, 21].

There is no strong direct limit over the b parameter, because of the difficulty of measuring a 2-Higgs state. However, an indirect limit arises because of the coupling between the hh decay and the elastic $\omega\omega$ scattering, as we will show later (see Ref. [68]). The direct limit over the a parameter, at a confidence level of 2σ ($\approx$95%), is

- CMS [70].. $a \in (0.87,\ 1.14)$
- ATLAS [69] .... $a \in (0.96,\ 1.34)$
- Fit of Buchalla et al. [71] .... $a \in (0.80,\ 1.16)$

The actual experimental results are shown in Fig. 2.2. Anyway, note that these limits are continuously improving, because of the LHC data reanalysis. Reference [71] uses the computer code Lilith-1.1.3 [72] for constraining the Effective Lagrangian. Giardino, in his Ph.D. Thesis [73], studied in detail the LHC constraints over a huge range of SM extensions, including the a parameter. We quote it though the situation is changing quickly.

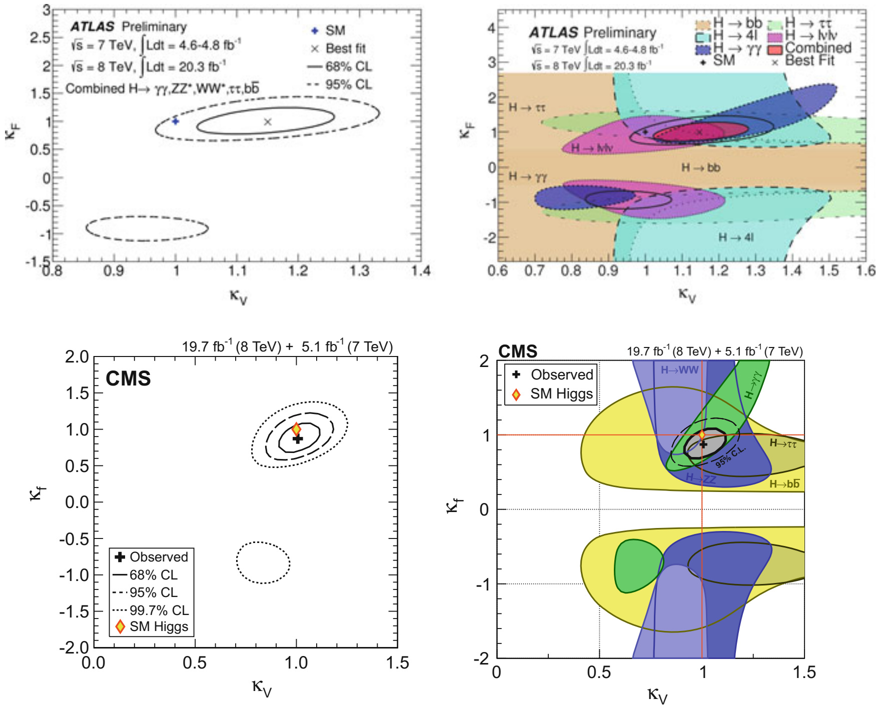

Fig. 2.2 From *top* to *bottom*, bounds over the $a = k_V$ and k_F coming from ATLAS ([69]; reproduced under the CC-BY-3.0 license) and CMS ([70]; reproduced under the CC-BY-4.0 license)

Now, let us consider the $U(\omega^a)$ coset. Since $U(\omega^a) \in SU(2)_L \times SU(2)_R / SU(2)_{L+R}$, it must be of the form

$$U = \quad + \frac{i\omega^a \tau^a}{v} + \mathcal{O}(\omega^2), \tag{2.12}$$

whatever the non-linear term is. So, the covariant derivative can be expanded as [67]

$$D_\mu U = \frac{i\partial_\mu \omega_a \tau^a}{v} + i\frac{g}{2} W_{\mu,i}\tau^i - i\frac{g'}{2} B_\mu \tau^3 + \cdots \tag{2.13}$$

However, specifying the parameterization of the $U \in SU(2)$ coset is necessary, since the non-linear terms will depend on it. One of the most usual elections is the exponential parameterization,

$$U(x) = \exp\left(i\frac{\tau^a \pi^a(x)}{v}\right), \tag{2.14}$$

τ^a $(a = 1, 2, 3)$ being the Pauli matrices.[5] Working with this expression,

$$U = \begin{pmatrix} \cos\left(\frac{\pi}{v}\right) + i\frac{\pi^3}{\pi}\sin\left(\frac{\pi}{v}\right) & \frac{i\pi^1+\pi^2}{\pi}\sin\left(\frac{\pi}{v}\right) \\ \frac{i\pi^1-\pi^2}{\pi}\sin\left(\frac{\pi}{v}\right) & \cos\left(\frac{\pi}{v}\right) - i\frac{\pi^3}{\pi}\sin\left(\frac{\pi}{v}\right) \end{pmatrix} = \cos\left(\frac{\pi}{v}\right) + i\frac{\tau^a\pi^a}{\pi}\sin\left(\frac{\pi}{v}\right), \tag{2.15}$$

where $\pi^2 \equiv \pi^a\pi^a$. With this parameterization, it can be checked that, as expected,

$$U \in SU(2) \Rightarrow U^\dagger \cdot U = U \cdot U^\dagger = \tag{2.16}$$

Only if we can neglect both the masses of Goldstone modes and couplings with longitudinal modes (ET limit, $M_W^2 \ll s$), then $D^\mu = \partial^\mu$ in Eq. 2.13. Otherwise, some of the couplings with $\hat{W}$ and $\hat{B}$ fields in Eq. 2.13 must we kept. Anyway,

$$\mathrm{Tr}\left[\partial_\mu U^\dagger \cdot \partial^\mu U\right] = \frac{A^2}{2v^2\pi^2} + \frac{\left(4v^2\pi^2 B - Av^2\right)}{2v^2\pi^4}\sin\frac{\pi}{v}, \tag{2.17}$$

where

$$\pi^2 = \pi^a\pi^a \tag{2.18a}$$
$$A = (\partial_\mu\pi^2)^2 = 4\left(\pi^a\pi^b\partial_\mu\pi^a\partial^\mu\pi^b\right) \tag{2.18b}$$
$$B = \partial_\mu\pi^a\partial^\mu\pi^a. \tag{2.18c}$$

Thus, $\mathrm{Tr}(\partial_\mu U^\dagger \cdot \partial^\mu U)$ can be expressed as

$$\mathrm{Tr}\left[\partial_\mu U^\dagger \cdot \partial^\mu U\right] = \frac{2}{v^2}\left[\frac{v^2}{\pi^2}\sin^2\frac{\pi}{v}\left(\delta_{ab} - \frac{\pi^a\pi^b}{\pi^2}\right) + \frac{\pi^a\pi^b}{\pi^2}\right]\partial_\mu\pi^a\partial^\mu\pi^b \tag{2.19}$$

However, as explained in [67], this is a good option for QCD Chiral Perturbation Theory (ChPT), where the $SU(3)$ coset is studied [41–44]. But we are dealing with an $SU(2)$ coset, which is isomorphic to S^3. Unless otherwise stated, the so-called *spherical parameterization* will be used (see Eq. 2.8), since computations in this basis are much simpler for the particular case of the $SU(2)$ coset. Note the notation change between Eqs. 2.8 and 2.14 ($\pi^a \leftrightarrow \omega^a$) in order to distinguish these two parameterizations. According to Eq. 2.15, we can change the parameterization by using

$$U = \cos\frac{\pi}{v} + i\frac{\tau^a\pi^a}{\pi}\sin\frac{\pi}{v} = \sqrt{1 - \frac{\omega^2}{v^2}} + i\frac{\tau^a\omega^a(x)}{v} \Rightarrow \frac{\omega^a}{v} = \frac{\pi^a}{\pi}\sin\frac{\pi}{v}. \tag{2.20}$$

By expanding this result,

$$\omega^a = \pi^a\left[1 - \frac{1}{6}\left(\frac{\pi}{v}\right)^2 + \frac{1}{120}\left(\frac{\pi}{v}\right)^4 - \frac{1}{5040}\left(\frac{\pi}{v}\right)^6 + \cdots\right]. \tag{2.21}$$

[5] Einstein's sum convention will be used unless otherwise stated.

Let us study the spherical basis. If Eq. 2.8 is expanded,

$$U = \begin{pmatrix} \sqrt{1-\frac{\omega^2}{v^2}} + \frac{i\omega^3}{v} & \frac{i\omega^1+\omega^2}{v} \\ \frac{i\omega^1-\omega^2}{v} & \sqrt{1-\frac{\omega^2}{v^2}} - \frac{i\omega^3}{v} \end{pmatrix}. \tag{2.22}$$

Equation 2.16 can be also checked within the spherical parameterization (Eq. 2.22). Now, computations are much simpler than with the exponential parameterization. For instance,

$$\mathrm{Tr}\left[\partial_\mu U_s^\dagger \cdot \partial^\mu U_s\right] = \frac{2}{v^2}\left[\delta_{ab} + \frac{\omega^a\omega^b}{v^2-\omega^2}\right]\partial_\mu\omega^a\partial^\mu\omega^b \tag{2.23}$$

For these fields ($\omega^a, \quad a = 1, 2, 3$) the spherical (or charge) basis is introduced,

$$\omega^\pm = \frac{\omega^1 \mp i\omega^2}{\sqrt{2}},\ \omega^0 = \omega^3, \tag{2.24}$$

which also implies

$$\omega^2 = 2\omega^+\omega^- + \omega^0\omega^0 \tag{2.25}$$

The same definition can be carried over to $\pi^\pm$ and π^0, which belongs to the exponential parameterization. On our Ref. [67], we studied these two parameterizations (spherical and exponential), for $\gamma\gamma$ scattering. As expected, the physical S-matrix elements are identical (in terms of ω^a and π^a, respectively). However, the intermediate results (i.e., the Feynman diagrams) are different.

According to the Equivalence Theorem (Sect. 2.1), if they are part of the physical initial or final states, and we have a high $\sqrt{s} \gg M_W$ (center of mass energy), the $\omega^\pm$, ω^0 (or $\pi^\pm$, π^0) can be identified with the longitudinal polarizations of gauge bosons $W^\pm$ and Z, up to an error $\mathcal{O}(M_W/\sqrt{s})$. The possible differences between parameterizations will be suppressed by a factor $\mathcal{O}(M_W/\sqrt{s})$. In the particular case of our Ref. [67], since we consider $M_W, M_h = 0$, these differences cancel.

2.2.1 WBGB Scattering

In Ref. [74] we reported the one-loop computation for the $\omega\omega \to \omega\omega$, $\omega\omega \to hh$ and $hh \to hh$ processes. Here we expand the discussion. Since we are working with an effective (non-renormalizable) theory, the following counterterms are needed in order to renormalize the scattering amplitudes,

$$\mathcal{L}_4 = a_4[\mathrm{Tr}(V_\mu V_\nu)][\mathrm{Tr}(V^\mu V^\nu)] + a_5[\mathrm{Tr}(V_\mu V^\mu)][\mathrm{Tr}(V_\nu V^\nu)] + \frac{\gamma}{f^4}(\partial_\mu\varphi\partial^\mu\varphi)^2 + \frac{\delta}{f^2}(\partial_\mu\varphi\partial^\mu\varphi)\mathrm{Tr}[(D_\nu U)^\dagger D^\nu U] + \frac{\eta}{f^2}(\partial_\mu\varphi\partial^\nu\varphi)\mathrm{Tr}[(D^\mu U)^\dagger D_\nu U] + \cdots, \tag{2.26}$$

where $V_\mu = (D_\mu U)U^\dagger$. The $\hat{W}_\mu U$ and $U\hat{B}_\mu$ terms of Eq. 2.13 are neglected, so that $D_\mu U \equiv \partial_\mu U$. This approximation is valid because we are in the Equivalence Theorem regime (Sect. 2.1) and we are dealing with couplings to neither photons nor transverse modes of gauge bosons. The spherical parameterization (see Eqs. 2.8 and 2.22) is used.

In subsequent work [68, 75] we changed the notation in Eq. 2.26, in order to adopt a recently agreed-upon standard. Thus, from now on, the next expressions will be used [67, 75]:

$$\mathcal{L}_0 = \frac{v^2}{4}\mathcal{F}(h)\mathrm{Tr}[(D_\mu U)^\dagger D^\mu U] + \frac{1}{2}\partial_\mu h\partial^\mu h - V(h) \tag{2.27}$$

$$\mathcal{L}_4 = a_4[\mathrm{Tr}(V_\mu V_\nu)][\mathrm{Tr}(V^\mu V^\nu)] + a_5[\mathrm{Tr}(V_\mu V^\mu)][\mathrm{Tr}(V_\nu V^\nu)] + \frac{g}{v^4}(\partial_\mu h\partial^\mu h)^2 + \frac{d}{v^2}(\partial_\mu h\partial^\mu h)\mathrm{Tr}[(D_\nu U)^\dagger D^\nu U] + \frac{e}{v^2}(\partial_\mu h\partial^\nu h)\mathrm{Tr}[(D^\mu U)^\dagger D_\nu U] + \cdots \tag{2.28}$$

$$\mathcal{F}(h) = 1 + 2a\frac{h}{v} + b\left(\frac{h}{v}\right)^2 + \cdots \tag{2.29}$$

$$V(h) = \sum_{n=0}^{\infty} V_n h^n = V_0 + \frac{1}{2}M_h^2 h^2 + d_3\frac{M_h^2}{2v}h^3 + d_4\frac{M_h^2}{8v^2}h^4 + \cdots \tag{2.30}$$

If we choose the spherical parameterization, according to [74, 75] and as follows from Eqs. 2.27 and 2.28, the NLO phenomenological Lagrangian for the WBGB scattering can be written as

$$\mathcal{L} = \frac{1}{2}\mathcal{F}(h)\partial_\mu\omega^a\partial^\mu\omega^b\left(\delta_{ab} + \frac{\omega^a\omega^b}{v^2}\right) + \frac{1}{2}\partial_\mu h\partial^\mu h + \frac{4a_4}{v^4}\partial_\mu\omega^a\partial_\nu\omega^a\partial^\mu\omega^b\partial^\nu\omega^b + \frac{4a_5}{v^4}\partial_\mu\omega^a\partial^\mu\omega^a\partial_\nu\omega^b\partial^\nu\omega^b + \frac{g}{v^4}(\partial_\mu h\partial^\nu h)^2 + \frac{2d}{v^4}\partial_\mu h\partial^\mu h\partial_\nu\omega^a\partial^\nu\omega^a + \frac{2e}{v^4}\partial_\mu h\partial^\nu h\partial^\mu\omega^a\partial_\nu\omega^a. \tag{2.31}$$

Note that we have neglected the Higgs mass and self-couplings which appear on the potential $V(h)$ that was defined in Eq. 2.30. This is valid on the regime $M_h \ll \sqrt{s}$, and provided that the strong dynamics is not triggered by unnaturally high d_i self-coupling parameters. There are direct experimental constraints over a_4 and a_5 parameters (Fig. 2.3). But not over g, d and e because of the difficulty of measuring a 2-Higgs state.

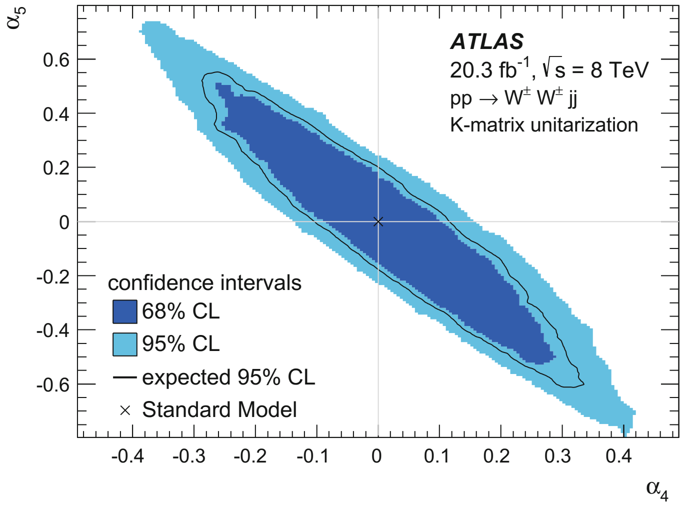

Fig. 2.3 Direct constraint over the a_4 and a_5 parameters coming from ATLAS [76]. The constraint of CMS [77] are given in terms of the F_{S0}/Λ^4 and F_{S1}/Λ^4 parameters, which have no direct translation to the a_4 and a_5 ones [78]. Figure reproduced from Ref. [76] under the CC-BY-3.0 license.

2.2.2 *Coupling with* $\gamma\gamma$

The effective Lagrangian (with the corresponding NLO counterterms) of Eq. 2.31 is valid provided that only interactions between WBGBs are taken into account. According to the Equivalence Theorem (Sect. 2.1), these WBGBs can be identified with the longitudinal modes of gauge bosons ($W^\pm$ and Z) and the Higgs-like scalar (H), as long as the CM energy is $\sqrt{s} \gg M_h^2, M_W^2, M_Z^2 \approx (100\,\text{GeV})^2$. Thus, for the processes $\gamma\gamma \to zz$ and $\gamma\gamma \to \omega^+\omega^-$, aditional terms involving the photon field must be introduced.

First of all, we can now not turn Eq. 2.13 into $D_\mu U = \partial_\mu U$ as in earlier works on the strong sector alone, because the couplings with the photon field A come from the couplings with $\hat{W}_{\mu\nu}$ and $\hat{B}_{\mu\nu}$ in Eq. 2.13, once a rotation to the physical basis is performed,

$$-c_W \frac{h}{v} \hat{W}_{\mu\nu} \hat{W}^{\mu\nu} - c_B \frac{h}{v} \hat{B}_{\mu\nu} \hat{B}^{\mu\nu} = -\frac{c_\gamma}{2} \frac{h}{v} e^2 A_{\mu\nu} A^{\mu\nu} + \cdots , \tag{2.32}$$

where $\hat{W}$ and $\hat{B}$ are defined in Eq. 2.10a and 2.10b. Thus, the $\mathcal{O}(p^2)$ Lagrangian [67] is, for the exponential parameterization,

$$\begin{aligned}\mathcal{L}_2(\pi, h, \gamma) = &\frac{1}{2}\partial_\mu h \partial^\mu h + \frac{1}{2}\mathcal{F}(h)(2\partial_\mu \pi^+ \partial^\mu \pi^- + \partial_\mu \pi^0 \partial^\mu \pi^0) \\ &+ \frac{1}{6v^2}\mathcal{F}(h)\left[(\partial_\mu \pi^+ \pi^- + \pi^+ \partial_\mu \pi^- + \pi^0 \partial_\mu \pi^0)^2 \right. \\ &\left. - \pi^2(2\partial_\mu \pi^+ \partial^\mu \pi^- + \partial_\mu \pi^0 \partial^\mu \pi^0)\right] \\ &+ ie\mathcal{F}(h)A^\mu(\partial_\mu \pi^+ \pi^- - \partial_\mu \pi^- \pi^+)\left(1 - \frac{\pi^2}{3v^2}\right) \\ &+ e^2\mathcal{F}(h)A_\mu A^\mu \pi^+ \pi^- \left(1 - \frac{\pi^2}{3v^2}\right). \end{aligned} \tag{2.33}$$

And, for the spherical one,

$$\begin{aligned}\mathcal{L}_2(\omega, h, \gamma) = &\frac{1}{2}\partial_\mu h \partial^\mu h + \frac{1}{2}\mathcal{F}(h)(2\partial_\mu \omega^+ \partial^\mu \omega^- + \partial_\mu \omega^0 \partial^\mu \omega^0) \\ &+ \frac{1}{2v^2}\mathcal{F}(h)(\partial_\mu \omega^+ \omega^- + \omega^+ \partial_\mu \omega^- + \omega^0 \partial_\mu \omega^0)^2 \\ &+ ie\mathcal{F}(h)A^\mu(\partial_\mu \omega^+ \omega^- - \omega^+ \partial_\mu \omega^-) + e^2\mathcal{F}(h)A_\mu A^\mu \omega^+ \omega^-. \end{aligned} \tag{2.34}$$

In both cases, $\mathcal{F}(h) = 1 + 2a(h/v) + b(h/v)^2$, as defined in Eq. 2.29. Furthermore, the next NLO extra counterterms are needed,

$$\mathcal{L}_{4'} = a_1 \mathrm{Tr}(U\hat{B}_{\mu\nu}U^\dagger \hat{W}^{\mu\nu}) + ia_2 \mathrm{Tr}(U\hat{B}_{\mu\nu}U^\dagger [V^\mu, V^\nu]) - ia_3 \mathrm{Tr}(\hat{W}_{\mu\nu}[V^\mu, V^\nu]), \tag{2.35}$$

besides those required for the WBGBs scattering: a_4, a_5, g, d and e (see Eq. 2.28). As for existing constraints on parameter values, Ref. [72] quotes a constraint on the c_γ parameter, coming from LHC data. At a confidence level of 2σ, $c_\gamma \in (-0.98, 0.50)/16\pi^2$. Note the $1/16\pi^2$ factor, which comes from the normalization of Ref. [72]. On the a_1, a_2 and a_3 parameters, there are only weaker constraints based on electroweak precision observables from LEP and purely SM one-loop calculations. See Ref. [79] for a review of such constraints. To sum up, $a_1 = (1.0 \pm 0.7) \times 10^{-3}$, $a_2 \in (-0.26, 0.26)$, $a_3 \in (-0.10, 0.04)$. However, these constraints are highly model dependent.

2.2.3 Coupling with $t\bar{t}$

The effective Lagrangian of Eq 2.31 models scattering processes between particles of the Electroweak Symmetry Breaking Sector. Thus, if we want to consider scattering processes between massive fermions, more terms (the so-called *Yukawa sector*) are required. Thus, as exposed in our Ref. [80], the considered effective Lagrangian is

$$\mathcal{L} = \frac{v^2}{4}\mathcal{F}(h)\mathrm{Tr}[(D_\mu U)^\dagger D^\mu U] + \frac{1}{2}\partial_\mu h \partial^\mu h - V(h) + i\bar{Q}\partial Q - v\mathcal{G}(h)[\bar{Q}'_L U H_Q Q'_R + h.c.]. \tag{2.36}$$

In the Yukawa sector of Eq. 2.36, the quark doublets are

$$Q' = \begin{pmatrix} \mathcal{U}' \\ \mathcal{D}' \end{pmatrix}, \tag{2.37}$$

where the two Q entries are made of the different up and down quark sectors

$$\mathcal{U}' = (u, c, t)', \qquad \mathcal{D}' = (d, s, b)'. \tag{2.38}$$

and the Yukawa-coupling matrix has the following form

$$H_Q = \begin{pmatrix} H_U & 0 \\ 0 & H_D \end{pmatrix}. \tag{2.39}$$

This matrix can be diagonalized by tranforming independently the right and left handed up and down quarks as:

$$\mathcal{D}_{L,R} = V^D_{L,R}\mathcal{D}'_{L,R}, \qquad \mathcal{U}_{L,R} = V^U_{L,R}\mathcal{U}'_{L,R}. \tag{2.40}$$

where $V^{U,D}_{L,R}$ are four 3×3 unitary matrices. Thus, the Yukawa part of the Lagrangian can be written as

$$\begin{aligned}\mathcal{L}_Y = -\mathcal{G}(h)\Bigg\{&\sqrt{1-\frac{\omega^2}{v^2}}\left(\bar{\mathcal{U}}M_U\mathcal{U} + \bar{\mathcal{D}}M_D\mathcal{D}\right) + \frac{i\omega^0}{v}\left(\bar{\mathcal{U}}M_U\gamma^5\mathcal{U} - \bar{\mathcal{D}}M_D\gamma^5\mathcal{D}\right)\\ &+\frac{i\sqrt{2}\omega^+}{v}\left(\bar{\mathcal{U}}_L V_{CKM}M_D\mathcal{D}_R - \bar{\mathcal{U}}_R M_U V_{CKM}\mathcal{D}_L\right)\\ &+\frac{i\sqrt{2}\omega^-}{v}\left(\bar{\mathcal{D}}_L V^\dagger_{CKM}M_U\mathcal{U}_R - \bar{\mathcal{D}}_R M_D V^\dagger_{CKM}\mathcal{U}_L\right)\Bigg\},\end{aligned} \tag{2.41}$$

where (Eq. 2.24) $\omega^\pm = (\omega^1 \mp i\omega^2)/\sqrt{2}$, $\omega^0 = \omega^3$ and the new quark fields are mass eigenstates, with M_U and M_D being the corresponding diagonal and real mass matrices. $V_{CKM} = V^U_L V^{D\dagger}_L$ is the Cabibbo-Kobayashi-Maskawa matrix. For the case of the heaviest quark generation (the only relevant for this work), the matrix element V_{tb} has been ommited since it is very close to unity. Hence, we have the effective Yukawa sector

$$\begin{aligned}\mathcal{L}_Y = -\mathcal{G}(h)\Bigg\{&\sqrt{1-\frac{\omega^2}{v^2}}\left(M_t t\bar{t} + M_b\bar{b}b\right) + \frac{i\omega^0}{v}\left(M_t\bar{t}\gamma^5 t - M_b\bar{b}\gamma^5 b\right)\\ &+\frac{i\sqrt{2}\omega^+}{v}\left(M_b\bar{t}_L b_R - M_t\bar{t}_R b_L\right) + \frac{i\sqrt{2}\omega^-}{v}\left(M_t\bar{b}_L t_R - M_b\bar{b}_R t_L\right)\Bigg\}.\end{aligned} \tag{2.42}$$

As can be seen, the different couplings of the left and right chiral parts of top and bottom quarks are an effect of custodial symmetry breaking. The $\mathcal{F}$ and $\mathcal{G}$ appearing in the Lagrangian (Eq. 2.36) are arbitrary analytical functions on the Higgs field h. $\mathcal{F}(h)$ is as defined in Eq. 2.29 above, $\mathcal{F}(h) = 1 + 2a(h/v) + b(h/v)^2$, whereas the new $\mathcal{G}(h)$ is

$$\mathcal{G}(h) = 1 + c_1 \frac{h}{v} + c_2 \frac{h^2}{v^2} + \cdots \tag{2.43}$$

For the computations presented in this work these functions are only needed up to the quartic terms $\mathcal{O}(h^4)$. Also we will consider the limit of massless botton quark ($M_b = 0$). Then we get the Yukawa Lagrangian:

$$\mathcal{L}_Y = -\mathcal{G}(h)\left[\left(1 - \frac{\omega^2}{2v^2}\right) M_t t\bar{t} + \frac{i\omega^0}{v} M_t \bar{t}\gamma^5 t - \frac{i\sqrt{2}\omega^+}{v} M_t \bar{t}_R b_L + \frac{i\sqrt{2}\omega^-}{v} M_t \bar{b}_L t_R\right], \tag{2.44}$$

where we have kept only the would-be Goldstone boson fields up to order $\omega^2 = 2\omega^+\omega^- + (\omega^0)^2$ (Eq. 2.25). Finally, the relevant Lagrangian for $\omega\omega \to t\bar{t}$ process in the regime[6] $M_h^2/v^2 \ll M_t^2/v^2 \ll s/v^2$ is given by

$$\begin{aligned}\mathcal{L} = -\mathcal{G}(h)&\left[\left(1 - \frac{\omega^2}{2v^2}\right) M_t t\bar{t} + \frac{i\sqrt{2}\omega^0}{v} M_t \bar{t}\gamma^5 t - \frac{i\sqrt{2}\omega^+}{v} M_t \bar{t}_R b_L + \frac{i\sqrt{2}\omega^-}{v} M_t \bar{b}_L t_R\right] \\ &+ \frac{1}{2}\mathcal{F}(h)\partial_\mu \omega^a \partial^\mu \omega_b \left(\delta_{ab} + \frac{\omega_a \omega_b}{v^2}\right) + \frac{1}{2}\partial_\mu h \partial^\mu h.\end{aligned} \tag{2.45}$$

Reference [71] also gives a direct constraint $c_1 \in (1.15, 1.53)$, at 2σ-confidence level. The one-loop divergences appearing in the $\omega\omega$ and hh scattering amplitudes can be absorbed in the GB four-derivative and the following two-derivative top antitop couplings

$$\begin{aligned}\mathcal{L}_4 = &\frac{4a_4}{v^4}\partial_\mu \omega^a \partial_\nu \omega^a \partial^\mu \omega^b \partial^\nu \omega^b + \frac{4a_5}{v^4}\partial_\mu \omega^a \partial^\mu \omega^a \partial_\nu \omega^b \partial^\nu \omega^b \\ &+ \frac{2d}{v^4}\partial_\mu h \partial^\mu h \partial_\nu \omega^a \partial^\nu \omega^a + \frac{2e}{v^4}\partial_\mu h \partial^\nu h \partial^\mu \omega^a \partial_\nu \omega^a + \frac{g}{v^4}\left(\partial_\mu h \partial^\mu h\right)^2 \\ &+ g_t \frac{M_t}{v^4}\partial_\mu \omega^a \partial^\mu \omega^b t\bar{t} + g_t' \frac{M_t}{v^4}\partial_\mu h \partial^\mu h t\bar{t}.\end{aligned} \tag{2.46}$$

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[6] This regime verifies the Equivalence Theorem. See Sect. 2.1.

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