Can the 125 GeV Higgs be the *Little Higgs*?

J. Reuter and M. Tonini

DESY Theory Group
Notkestr. 85, 22603 Hamburg, Germany

E-mail: juergen.reuter@desy.de, marco.tonini@desy.de

ABSTRACT: After the discovery of the Higgs-like boson by the LHC 2012 it is the most important task to check whether this new particle is the Standard Model Higgs boson or something else. In this paper, we study whether the 125 GeV boson could be the pseudo-Goldstone boson of Little Higgs models. We derive limits on the parameter space of several Little Higgs models (simple group and product group models, with and without T-parity), both from the experimental data from ATLAS and CMS about the different Higgs discovery channel and the electroweak precision observables. We perform a fit of several Little Higgs models to all electroweak parameters from measurements of SLC, LEP, Tevatron, and LHC. For the Higgs searches, we include all available data from the summer conferences in 2012 as well as the updates from December 2012. We show that there always exists a region in the parameter space of the models under consideration where the measured χ^2 is equal or lower than the SM χ^2: a closer look at the minimum χ^2 will however reveal that the agreement with the collected data is not significantly better as within the SM. While for the models without T-parity the Little Higgs scale f is forced to be of the order 2-4 TeV in order to be compatible with the collected data, in the models with T-parity the scale f is constrained to be only above $\mathcal{O}(500)$ GeV, reducing the amount of fine-tuning. We also show that these results are still driven by the electroweak precision measurements due to the bigger LHC data uncertainties.

Contents

1 Introduction

The discovery of a bosonic particle with a mass of 125 GeV by the LHC experiments
2012 [1, 2] seems to be the last piece of the jigsaw puzzle of the electroweak interactions.
However, at present it is not yet clear whether this particle does really have all the properties
of the Standard Model (SM) Higgs boson, or whether it is a particle of some extension of
the SM.

There are many reasons to believe in the existence of beyond the Standard Model (BSM)
physics: the missing CP violation needed for the explanation of the baryon-antibaryon
asymmetry in the universe, the missing dark matter component in the SM, and the question
about the stability of the Higgs potential and the electroweak vacuum: the latter has been
called the hierarchy or fine-tuning problem, namely the problem that the Higgs self-coupling
is driven to non-perturbative values for too large Higgs masses while the top couplings tend
to destabilize the electroweak vacuum. Furthermore, the bare Higgs mass parameter seems
to be tuned very accurately in order to get a Higgs mass at the electroweak scale, as scalar
masses are quadratically sensitive to new physics particles coupling to them.

One paradigm to solve this problem is to assume the Higgs boson to be no fundamental,
but a composite particle, as e.g. in Technicolor, Topcolor or composite Higgs models.
The Higgs boson is relatively light compared to high scales, because it appears, like the
pions in chiral symmetry breaking, as the (pseudo)-Nambu-Goldstone bosons (pNGBs)

of a spontaneously broken global symmetry. However, this necessitates the presence of strong interactions to bind new constituents together to something like the Higgs boson, and indications of such strong interactions had to show up in the electroweak precision measurements from SLC and LEP (and also Tevatron). A solution to this problem has been found using the formalism of collective symmetry breaking, where several global symmetries are intertwined. If each of them were exact, the Higgs would still be an exact massless Goldstone boson. Hence, the mass term arises only logarithmically at the one-loop order or quadratically at two-loop order. This leaves such models weakly interacting at the TeV scale and raises the scale for the onset of new strong interactions to several TeV up to tens of TeV. These kinds of models have first been realized motivated from deconstructed extra dimensions [3, 4], and then in a 4D setup by explicit constructions of coset spaces for the symmetry breaking pattern [5–7]. There are two different types of models, so-called Simple Group Models, where the weak gauge group extension is given by a simple Lie group, while the Goldstone multiplet of the broken global symmetry is distributed over several different non-linear sigma model fields, whereas in the Product Group Models the Goldstone multiplet is a single representation parameterizing the coset space of the global symmetry breaking, and the weak gauge group emerges as the unbroken part of a product gauge group. The most prominent examples of these two classes are the Simplest Little Higgs model [8] and the Littlest Higgs model [5], respectively. To ameliorate the amount of fine tuning within the so-called Little Hierarchy problem between the electroweak and the TeV scale, a discrete symmetry named T-parity has been introduced [9]: it cancels tree-level contributions from heavy Little Higgs states to the electroweak precision observables (at least in the gauge and scalar sector for the product-group models), and offers a possibility for a dark matter particle. For an overview over and more details about Little Higgs models, cf. [10, 11].

In this paper, we discuss the most common Little Higgs models, namely the Littlest Higgs with and without T-parity as well as the Simplest Little Higgs, and fit the results reported by both experimental collaborations, ATLAS and CMS, about the many different Higgs search/discovery channels to these different models. This is accompanied by a simultaneous fit of the electroweak precision data to these models. We compare the constraints coming from the LHC Higgs discovery with those from electroweak precision physics. The paper is organized as follows: Sec. 2 gives a technical introduction into the three different Little Higgs models under consideration, structured according to their gauge and scalar sector, the fermion sector, and finally discussing the electroweak precision observables for these models. The experimental data needed for the analysis presented here are given in Sec. 3, where we also present the statistical methods that we used to perform the fit of the Little Higgs models to the experimental data as well as to the precision observables. Our results are presented in Sec. 4, before we give our conclusions in Sec. 5 . In the appendix, technical details on the determination of the Higgs boson partial widths and cross section as well as on the calculation of the electroweak precision observables within the Little Higgs models are shown.

2 The Little Higgs framework

In this section, we will describe the structure of the three different Little Higgs models under consideration, focusing in particular on the details which will affect our results. However, this section should not be thought as a comprehensive review of these models, for which we refer to [10, 11].

We decided to present separately the structure of the gauge, the scalar and the fermion sectors, in order to underline the different implementations of the Little Higgs paradigm in the considered models. In the end, there is also a subsection describing the effect of the Little Higgs structure on the predictions of Electroweak Precision Observables (EWPO).

2.1 Gauge and Scalar sectors

Littlest Higgs Model

The Littlest Higgs model (we mainly follow the presentation given in [12] instead of the original paper [5]; in the sequel, we use the abbreviation L^2H) is based on a non-linear sigma model in the coset space

$$SU(5)/SO(5). \tag{2.1}$$

The vacuum expectation value (*vev*) of an $SU(5)$ symmetric tensor field generates the global spontaneous symmetry breaking (2.1) at the scale f:

$$\langle \Sigma \rangle = \begin{pmatrix} \mathbf{0}_{2\times2} & \mathbf{0}_{2\times1} & \mathbf{1}_2 \\ \mathbf{0}_{1\times2} & 1 & \mathbf{0}_{1\times2} \\ \mathbf{1}_2 & \mathbf{0}_{2\times1} & \mathbf{0}_{2\times2} \end{pmatrix} \tag{2.2}$$

In this setup, there are 14 Nambu-Goldstone Bosons (NGBs) Π^a, $a = 1, \ldots, 14$, parametrized by

$$\Sigma(x) = e^{2i\Pi^a X^a(x)/f} \langle \Sigma \rangle \tag{2.3}$$

where X^a are the broken generators of the coset space $SU(5)/SO(5)$.

This model belongs to the class of Product Group models, where the SM gauge group emerges from the diagonal breaking of the product of several gauged groups: in this specific realization there is a local invariance under $[SU(2)_1 \otimes U(1)_1] \otimes [SU(2)_2 \otimes U(1)_2]$, embedded in the matrix structure, spontaneously broken through the *vev* $\langle \Sigma \rangle$ to its diagonal subgroup, which is identified with the SM gauge group. A set of $SU(2) \otimes U(1)$ gauge bosons obtains a mass of order f, while the other set is left massless and is identified with the SM gauge fields.

Under the unbroken $SU(2)_L \otimes U(1)_Y$ the Π^a transform as $\mathbf{1}_0 \oplus \mathbf{3}_0 \oplus \mathbf{2}_{1/2} \oplus \mathbf{3}_{\pm1}$: the $\mathbf{2}_{1/2}$ component is identified with the Higgs boson h, while the $\mathbf{3}_{\pm1}$ component is a complex triplet under $SU(2)_L$ which forms a symmetric tensor $\Phi_{ij} \equiv \Phi$ with components ϕ^{++}, ϕ^+, ϕ^0 and a pseudo-scalar ϕ^P, where both ϕ^0 and ϕ^P are real scalars. The other components are the longitudinal modes of the heavy gauge bosons and therefore will not appear in unitary gauge.

The kinetic term for the NGB matrix can be expressed in the standard non-linear sigma model formalism as

$$\mathcal{L}_\Sigma = \frac{1}{2}\frac{f^2}{4}\,\mathrm{tr}\big|D_\mu\Sigma\big|^2 \tag{2.4}$$

where the numerical coefficients assure canonically normalized kinetic terms for the scalar fields. To impose a local invariance under $[SU(2)_1 \otimes U(1)_1] \otimes [SU(2)_2 \otimes U(1)_2]$, the covariant derivative is defined as

$$D_\mu\Sigma = \partial_\mu\Sigma - i\sum_{j=1}^{2}\big[g_j(W_j\Sigma + \Sigma W_j^t) + g_j'(B_j\Sigma + \Sigma B_j^t)\big] \tag{2.5}$$

and the generators of the gauged symmetries are explicitly given as

$$Q_1^a = \begin{pmatrix} \sigma^a/2 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix} \qquad Y_1 = \mathrm{diag}\,(3,3,-2,-2,-2)/10$$

$$Q_2^a = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & -\sigma^{a*}/2 \end{pmatrix} \qquad Y_2 = \mathrm{diag}\,(2,2,2,-3,-3)/10. \tag{2.6}$$

The global symmetries prevent the appearance of a potential for the scalar fields at tree level. The scalar potential is indeed generated at one-loop and higher orders due to the interactions with gauge bosons and fermions, and is parametrized through the Coleman-Weinberg (CW) potential [13]. The scalar potential takes the generic form

$$V_{CW} = \lambda_{\phi^2}f^2\,\mathrm{tr}(\phi^\dagger\phi) + i\lambda_{h\phi h}f(h\phi^\dagger h^T - h^*\phi h^\dagger) - \mu^2 hh^\dagger + \lambda_{h^4}(hh^\dagger)^2 \tag{2.7}$$

where the coefficients λ_{ϕ^2}, $\lambda_{h\phi h}$ and λ_{h^4} are functions of the fundamental parameters of the model, while the Higgs mass parameter μ^2 should be treated as a free parameter since it receives big contributions also from two-loop diagrams, that have not been calculated.

Minimizing the potential to obtain the doublet and triplet *vev*s v and v', and requiring appropriate relations to correctly trigger electroweak symmetry breaking (EWSB), one can express all four parameters in the scalar potential to leading order in terms of the physical parameters f, m_h^2, v and v', and obtain the following relation between the two *vev*s

$$x \equiv \frac{4v'f}{v^2}, \qquad 0 \le x < 1\,. \tag{2.8}$$

Diagonalizing the scalar mass matrix, one obtains at leading order the following spectrum:

$$m_h = \sqrt{2}\mu, \qquad m_\Phi = \frac{\sqrt{2}m_h}{\sqrt{1-x^2}}\frac{f}{v} \qquad , \tag{2.9}$$

where all components of the triplet $\big(\phi^{++}, \phi^+, \phi^0, \phi^P\big)$ are degenerate at the order we are considering. Since μ^2 is treated as a free parameter, we will assume the measured Higgs mass for the scalar doublet h, fixing therefore the value of μ.

If we parametrize the interaction terms of the charged components of the triplet to the Higgs field in the following way

$$V_{CW} \supset -2\,\frac{m_\Phi^2}{v}\,y_{\phi^+}\,\phi^+\phi^-h - 2\,\frac{m_\Phi^2}{v}\,y_{\phi^{++}}\,\phi^{++}\phi^{--}h \tag{2.10}$$

then the couplings y_ϕ after EWSB are predicted up to $\mathcal{O}\left(v^2/f^2\right)$ to be [14]

$$y_{\phi^+} = \frac{v^2}{f^2}\left(-\frac{1}{3} + \frac{x^2}{4}\right), \qquad y_{\phi^{++}} = \mathcal{O}\left(\frac{v^4}{f^4}\right). \tag{2.11}$$

This parametrization will be useful to calculate the contribution of the charged resonances to the one-loop $h\gamma\gamma$ vertex, cf. Appendix A.

A set of $SU(2) \otimes U(1)$ gauge bosons (W', B') obtains a mass term of order f from (2.4), while the other set (W, B) remains massless. The mass eigenstates are related to the gauge eigenstates by the following field rotations

$$W = sW_1 + cW_2, \qquad W' = -cW_1 + sW_2 \tag{2.12}$$
$$B = s'B_1 + c'B_2 \qquad B' = -c'B_1 + s'B_2$$

where the mixing angles, which we will treat as free parameters, are given by

$$c = \frac{g_1}{\sqrt{g_1^2 + g_2^2}}, \qquad c' = \frac{g_1'}{\sqrt{g_1'^2 + g_2'^2}}. \tag{2.13}$$

EWSB induces further mixing between the light and heavy gauge bosons: at leading order, the spectrum is given by

$$\begin{aligned} m_{W^\pm} &= \frac{gv}{2} & m_{W_H^\pm} &= \frac{gf}{2sc} \\ m_Z &= \frac{gv}{2c_w} & m_{Z_H} &= \frac{gf}{2sc} \\ m_\gamma &= 0 & m_{A_H} &= \frac{g'f}{2\sqrt{5}s'c'} \end{aligned} \tag{2.14}$$

If we parametrize the interaction terms of the charged gauge bosons to the Higgs field as

$$\mathcal{L}_\Sigma \supset 2\,\frac{m_W^2}{v}\,y_W\,W^+W^-h + 2\,\frac{m_{W_H}^2}{v}\,y_{W_H}\,W_H^+W_H^-h \tag{2.15}$$

then the couplings y_V after EWSB are predicted up to $\mathcal{O}\left(v^2/f^2\right)$ to be [14]

$$y_W = 1 + \frac{v^2}{f^2}\left[-\frac{1}{6} - \frac{1}{4}(c^2 - s^2)^2\right], \qquad y_{W_H} = -s^2c^2\frac{v^2}{f^2}. \tag{2.16}$$

The L^2H model contains new matter content and interactions which contribute to the EWPO, as we will discuss in detail later. In particular, from the exchange of heavy $SU(2)$ gauge bosons and from the presence of the triplet *vev* v', the relation between the Fermi constant G_F and the doublet *vev* v is modified from its SM form: by comparing the two

relations one can thus express the L^2H doublet *vev* v in terms of the SM value $v_{SM} = 246$ GeV up to $\mathcal{O}\left(v_{SM}^2/f^2\right)$ as [14]

$$v = v_{SM}\left[1 - \frac{v_{SM}^2}{f^2}\left(-\frac{5}{24} + \frac{x^2}{8}\right)\right]. \tag{2.17}$$

Using this relation, we can express the corrections of the SM-like hVV couplings ($V \equiv W$, Z) with respect to their SM values up to $\mathcal{O}\left(v_{SM}^2/f^2\right)$ equivalently as

$$\frac{g_{hVV}}{g_{hVV}^{SM}} = 1 + \frac{1}{8}\frac{v_{SM}^2}{f^2}\left[-3 + x^2 - 2(c^2 - s^2)^2\right] \tag{2.18}$$

where

$$g_{hVV} = \frac{m_V^2}{v}y_V, \qquad g_{hVV}^{SM} = \frac{m_V^2}{v}\bigg|_{v=v_{SM}}. \tag{2.19}$$

Eq. (2.18) will be useful to calculate the tree-level decays of the Higgs boson into the SM-like gauge bosons, cf. Appendix A .

We will not consider all other tree-level decay channels of the Higgs which involve the heavy gauge bosons or the heavy scalar triplet: indeed in L^2H the EWPD require f larger than a few TeV, cf. Ref. [15] and our results of Sec. 4, making these decay channels kinematically forbidden.

Littlest Higgs with T-parity

As just mentioned, the original Littlest Higgs model suffers from severe constraints from EWPO, which could only be satisfied in small regions of the parameter space. The most severe constraints resulted from tree-level corrections to EWPO due to the exchange of the heavy gauge bosons present in the theory, as well as from the small but non-vanishing *vev* of the additional scalar triplet field Φ. These severe constraints are evaded with the introduction of a conserved discrete symmetry, called *T-parity*, featuring T-odd partners for all (T-even) SM particles, and a lightest T-odd particle that is stable. As a result, tree-level contributions of the heavy gauge bosons to EWPO are suppressed, and corrections arise only at loop level.

The Littlest Higgs model with T-parity (for detailed reviews cf. [17, 18], and the original papers [9, 19]; in the following we use the abbreviation LHT) shares the same global and local symmetry structure of the original L^2H model. The LHT model has therefore the same scalar kinetic term of Eq. (2.4), where the T-parity can be naturally implemented requiring that the coupling constant of $SU(2)_1$ ($U(1)_1$) equals that of $SU(2)_2$ ($U(1)_2$): in this way the four mixing angles of the gauge sector c, s, c', s' are all equal to $1/\sqrt{2}$.

Under T-parity, the Higgs field and the SM-like gauge bosons are T-even, while the scalar triplet and the heavy gauge bosons are T-odd. Therefore the coupling $h^\dagger\Phi h$ is forbidden, leading to the relations for the triplet *vev* $v' = 0$ and $x = 0$. Since the correction of W_H to the relation between G_F and v is forbidden by T-parity, the functional form of the Higgs *vev* v up to $\mathcal{O}\left(v_{SM}^2/f^2\right)$ is modified as [14]

$$v = v_{SM}\left(1 + \frac{1}{12}\frac{v_{SM}^2}{f^2}\right). \tag{2.20}$$

The scalar and gauge boson mass spectrum and their couplings to the Higgs field in *LHT* model can be easily obtained from the respective L^2H relations by taking $c = s = c' = s' = 1/\sqrt{2}$ and $x = 0$. Only the hVV coupling ($V \equiv W, Z$) gets a different correction in the *LHT* model because of the different functional form of v [20]:

$$\frac{g_{hVV}}{g_{hVV}^{SM}} = 1 - \frac{1}{4}\frac{v_{SM}^2}{f^2} - \frac{1}{32}\frac{v_{SM}^4}{f^4} + \mathcal{O}\left(\frac{v_{SM}^6}{f^6}\right). \tag{2.21}$$

If the lightest T-odd particle A_H is very light, also the tree-level decay $h \to A_H A_H$ could be kinematically open in *LHT*. The $hA_H A_H$ coupling is given by [12]

$$g_{hA_H A_H} = -\frac{1}{2}g'^2 v \qquad , \tag{2.22}$$

and in Appendix A there is the explicit expression of the partial width of this decay channel: indeed in *LHT* a lower value of f is allowed by EWPO [18], and thus this decay channel could be kinematically open. Note that if one assumes the A_H to be the dark matter particle, a resonant coannihilation of two heavy photons via s-channel Higgs exchange is actually favored, rendering this channel close to irrelevant.

Simplest Little Higgs

The Simplest Little Higgs (for detailed reviews cf. [21–23], while the original references are [7, 8], and the used abbreviation *SLH*) is based on a non-linear sigma model in the coset space

$$\frac{[SU(3)_1 \otimes U(1)_1] \otimes [SU(3)_2 \otimes U(1)_2]}{[SU(2)_1 \otimes U(1)_1] \otimes [SU(2)_2 \otimes U(1)_2]}. \tag{2.23}$$

The *vev*s of two $SU(3)_1 \otimes SU(3)_2$ scalar fields $\phi_1 \sim (\mathbf{3}, \mathbf{1})$ and $\phi_2 \sim (\mathbf{1}, \mathbf{3})$ realize the spontaneous symmetry breaking $SU(3)_i \to SU(2)_i$ ($i = 1, 2$) at scales f_1 and f_2 respectively, giving rise to ten NGBs.

This model belongs to the class of Simple Group models, where the SM gauge group emerges from the breaking of a larger simple group: in this specific realization there is a local invariance under the diagonal subgroup $SU(3)_L \otimes U(1)_X$, which is spontaneously broken by the *vev*s of $\phi_{1,2}$ to the SM $SU(2)_L \otimes U(1)_Y$. Five NGBs are therefore eaten and five new gauge bosons arise with a mass of the order of the scale f, with $f^2 = f_1^2 + f_2^2$.

The NGBs are parametrized with a non-linear representation of the two complex scalar triplet fields $\phi_{1,2}$

$$\phi_1(x) = \exp\left(\frac{it_\beta \Theta(x)}{f}\right)\begin{pmatrix} 0 \\ 0 \\ fc_\beta \end{pmatrix}, \qquad \phi_2(x) = \exp\left(-\frac{i\Theta(x)}{t_\beta f}\right)\begin{pmatrix} 0 \\ 0 \\ fs_\beta \end{pmatrix} \tag{2.24}$$

with $t_\beta = \sin\beta/\cos\beta = f_2/f_1$ being the ratio of the *vev*s of the scalar triplets, and $\Theta(x)$ the NGB matrix

$$\Theta = \frac{1}{f}\left[\begin{pmatrix} \mathbf{0}_{2\times 2} & h \\ h^\dagger & 0 \end{pmatrix} + \frac{\eta}{\sqrt{2}}\mathbf{1}_{3\times 3}\right]. \tag{2.25}$$

Here, we have already neglected (in unitary gauge) the NGBs that become the longitudinal modes of the 5 heavy and 3 SM-like gauge bosons (after EWSB): indeed the remaining 2 physical NGBs are identified with the Higgs doublet h and with a pseudo-scalar η as above.

The presence of the pseudo-scalar η and in particular of the coupling h-Z-η is a peculiar and distinguishing feature of Simple Group models class, as already pointed out in [24].

The kinetic term for the scalar sector can be expressed in the standard non-linear sigma model formalism as

$$\mathcal{L}_\Phi = \sum_{i=1}^{2} \left| D_\mu \phi_i \right|^2 \tag{2.26}$$

where the covariant derivative, in order to assure $SU(3)_L \otimes U(1)_X$ local invariance, is

$$D_\mu = \partial_\mu - igA_\mu^a T_a + ig_x Q_x B_\mu^x, \qquad g_x = \frac{g'}{\sqrt{1 - t_w^2/3}} \tag{2.27}$$

with $t_W \equiv \tan\theta_W$, and g, g' the SM $SU(2)_L \otimes U(1)_Y$ gauge couplings.

The global symmetries prevent the appearance of a Higgs potential at tree level. The Higgs potential is indeed generated at one-loop and higher orders due to the interactions with gauge bosons and fermions through the CW potential. One can show [22] that in this setup the pseudo-scalar η remains massless, while the Higgs boson acquires a mass through one-loop logarithmic and two-loop quadratic divergences (this is due to the collective symmetry breaking mechanism). To force η to have a non-zero mass, in order to avoid a new and not observed long-range interaction, one possible solution is to introduce a term

$$-\mu_\phi^2 \left(\phi_1^\dagger \phi_2 + \text{h.c.} \right) \tag{2.28}$$

into the CW potential by hand. This explicitly breaks the global $SU(3)$ symmetry and also the collective symmetry breaking mechanism, but the corrections are small [22]: we will adopt this extension, and the parameter μ_ϕ will be then proportional to the pseudo-scalar mass m_η. The CW potential now becomes

$$V_{CW} = -\mu^2 h^\dagger h + \lambda (h^\dagger h)^2 - \frac{1}{2} m_\eta^2 \eta^2 + \lambda' h^\dagger h \eta^2 + \dots \tag{2.29}$$

where the parameters are defined as in [22]. From the minimization of the potential one obtains the expression for the *vev* of the Higgs field

$$v^2 = \frac{\mu^2}{\lambda} \tag{2.30}$$

and the mass of the pseudo-scalar η

$$m_\eta^2 = \frac{\mu_\phi^2}{c_\beta s_\beta} \cos\left(\frac{v}{\sqrt{2} f s_\beta c_\beta} \right). \tag{2.31}$$

If one assumes that the physics at the cut-off $\Lambda = 4\pi f$ gives no sizable contribution to the scalar potential, then the CW potential (2.29) fully determines the scalar masses and their couplings: in order to realize a correct EWSB pattern, the free parameters of the CW

potential are then not anymore independent among themselves. In particular, we require the parameter μ to reproduce the observed Higgs boson mass

$$m_h = \sqrt{2}\,\mu \tag{2.32}$$

while we fix v in Eq. (2.30) in order to match the prediction of the SM W-boson mass: the W-boson mass is indeed predicted to be [22]

$$m_W = \frac{gv}{2}\left[1 - \frac{1}{12}\frac{v^2}{f^2}\frac{t_\beta^4 - t_\beta^2 + 1}{t_\beta^2} + \frac{1}{180}\frac{v^4}{f^4}\frac{t_\beta^8 - t_\beta^6 + t_\beta^4 - t_\beta^2 + 1}{t_\beta^4} + \mathcal{O}\left(\frac{v^6}{f^6}\right)\right] \tag{2.33}$$

and therefore we require v to satisfy

$$v \simeq v_{SM}\left[1 + \frac{1}{12}\frac{v_{SM}^2}{f^2}\frac{t_\beta^4 - t_\beta^2 + 1}{t_\beta^2} - \frac{1}{180}\frac{v_{SM}^4}{f^4}\frac{t_\beta^8 - t_\beta^6 + t_\beta^4 - t_\beta^2 + 1}{t_\beta^4}\right] \tag{2.34}$$

$$\equiv v_{SM}\left[1 + \delta_v^{(2)} - \delta_v^{(4)}\right]$$

where $v_{SM} = 246$ GeV, in order to have $m_W = g\,v_{SM}/2$. Therefore (2.32) and (2.30, 2.34) are two conditions which have to be imposed on the free parameters of the CW potential, i.e. on f, t_β, μ_ϕ, R (R is a ratio of Yukawa couplings of the fermion sector which affects μ, λ, cf. below): we decided to let f and t_β to be free parameters of our study, fixing the values of μ_ϕ and R through the previous equations.

From (2.34) we also see that the correction to v_{SM} is proportional to $t_\beta^2 v_{SM}^2/f^2$ in the large t_β limit: as suggested in [22], for perturbation theory to be valid, the $\mathcal{O}\left(v_{SM}^4/f^4\right)$ correction should be suppressed by a factor of 0.1 relative to the $\mathcal{O}\left(v_{SM}^2/f^2\right)$ correction, i.e.

$$\delta_v^{(4)}/\delta_v^{(2)} < 0.1. \tag{2.35}$$

We will require this latter condition to be satisfied in the considered parameter space of the model.

After EWSB and using relation (2.34), the leading order mass spectrum of the heavy and light (SM) gauge bosons is given by [21, 22]

$$m_{W^\pm} = \frac{g\,v_{SM}}{2} \qquad m_{X^\pm} = m_{Y^0} = m_{\bar{Y}^0} = \frac{gf}{\sqrt{2}}$$

$$m_Z = \frac{g\,v_{SM}}{2c_w}\left(1 + \frac{v^2}{16f^2}(1 - t_w^2)^2\right) \qquad m_{Z'} = \sqrt{\frac{2}{3 - t_w^2}}gf \tag{2.36}$$

$$m_\gamma = 0$$

where we have included also the $\mathcal{O}(v^2/f^2)$ custodial symmetry violating shift term in the Z-mass, and c_w is the cosine of the Weinberg angle. If we parametrize the interaction terms of the charged gauge bosons to the Higgs field in the following way

$$\mathcal{L}_\Phi \supset 2\frac{m_W^2}{v}y_W\,W^+W^-h + 2\frac{m_X^2}{v}y_X\,X^+X^-h \tag{2.37}$$

then the couplings y_V after EWSB are predicted to be [14]

$$y_W \simeq \frac{v}{v_{SM}} \left[1 - \frac{1}{4} \frac{v_{SM}^2}{f^2} \frac{t_\beta^4 - t_\beta^2 + 1}{t_\beta^2} + \frac{1}{36} \frac{v_{SM}^4}{f^4} \frac{\left(t_\beta^2 - 1\right)^2}{t_\beta^2} \right], \qquad y_X \simeq -\frac{1}{2} \frac{v^2}{f^2} . \qquad (2.38)$$

Using relation (2.34), we can express the corrections of the hVV ($V \equiv W, Z$) couplings with respect to their SM value up to $\mathcal{O}\left(v_{SM}^4/f^4\right)$ equivalently as [22]

$$\frac{g_{hWW}}{g_{hWW}^{SM}} = 1 - \frac{1}{4} \frac{v_{SM}^2}{f^2} \left(\frac{t_\beta^4 - t_\beta^2 + 1}{t_\beta^2} \right) + \frac{1}{36} \frac{v_{SM}^4}{f^4} \frac{\left(t_\beta^2 - 1\right)^2}{t_\beta^2} \qquad (2.39)$$

$$\frac{g_{hZZ}}{g_{hZZ}^{SM}} = 1 - \frac{1}{4} \frac{v_{SM}^2}{f^2} \left(\frac{t_\beta^4 - t_\beta^2 + 1}{t_\beta^2} + \left(1 - t_w^2\right)^2 \right) + \frac{1}{36} \frac{v_{SM}^4}{f^4} \frac{\left(t_\beta^2 - 1\right)^2}{t_\beta^2}$$

where as usual

$$g_{hVV} = \frac{m_V^2}{v} y_V, \qquad g_{hVV}^{SM} = \frac{m_V^2}{v} y_V \Big|_{v=v_{SM}} . \qquad (2.40)$$

2.2 Fermion sector

Littlest Higgs

The SM fermions acquire their masses through the Higgs mechanism via Yukawa interactions: the large top Yukawa coupling induces a dominant quadratic correction to the Higgs boson mass, spoiling the naturalness of a light Higgs boson. In L^2H model this problem is solved by introducing a new set of heavy fermions with coupling to the Higgs field such that it cancels the quadratic divergence due to the top quark. The new fermions are a vectorlike pair (T', T'^c) with quantum numbers $(\mathbf{3}, \mathbf{1})_{Y_i}$, $(\bar{\mathbf{3}}, \mathbf{1})_{-Y_i}$ respectively, and therefore they are allowed to have a bare mass term which is chosen to be of order f.

The Yukawa-like Lagrangian for the third generation of quarks can be found e.g. in [12], and contains the following interaction terms to the Higgs after EWSB:

$$\mathcal{L}_t \supset -\lambda_1 f \left(\frac{s_\Sigma}{\sqrt{2}} \bar{t}_L t'_R + \frac{1 + c_\Sigma}{2} \bar{T}'_L t'_R \right) - \lambda_2 f \bar{T}'_L T'_R + \text{h.c.} \qquad (2.41)$$

where $c_\Sigma = \cos\left(\sqrt{2}h/f\right)$, $s_\Sigma = \sin\left(\sqrt{2}h/f\right)$, and with $\lambda_{1,2}$ as free parameters. After diagonalization of the mass matrix, the leading order mass eigenvalues are the following

$$m_t = \frac{\lambda_2 R}{\sqrt{1 + R^2}} v, \qquad m_T = \lambda_2 \sqrt{1 + R^2} f . \qquad (2.42)$$

Here, we have defined the ratio of the Yukawa couplings

$$R = \lambda_1/\lambda_2. \qquad (2.43)$$

However we can fix λ_2 requiring that, for given (f, R), m_t corresponds to the experimental top mass value: in this way, the only free parameters in the top sector are f and R. If we

parametrize the interaction terms of the top quark and heavy top to the Higgs field (the dominant contributions to the effective hgg vertex) in the following way

$$\mathcal{L}_t \supset -\frac{m_t}{v}\, y_t\, \bar{t}\, t\, h - \frac{m_T}{v}\, y_T\, \bar{T}\, T\, h \qquad , \tag{2.44}$$

then the couplings $y_{t,T}$ after EWSB are predicted up to $\mathcal{O}\left(v^2/f^2\right)$ to be [14]

$$y_t = 1 - \frac{v^2}{f^2}\left[\frac{2R^4 + R^2 + 2}{3(1+R^2)^2} + \frac{x^2}{4} - \frac{x}{2}\right], \qquad y_T = -\frac{R^2}{(1+R^2)^2}\frac{v^2}{f^2}. \tag{2.45}$$

The scalar interactions with the *up*-type quarks of the first two generations have the same form as $\mathcal{L}_t$, except that there is no need for extra vectorlike quarks. The interactions with the *down*-type quarks and leptons of the three generations are generated by a similar Lagrangian, again without the extra vectorlike quarks. For the explicit forms of the Lagrangian terms we refer as before to Ref. [12]. The important result for our analysis is the explicit correction of the Higgs-fermion couplings with respect to their SM value: from the Feynman rules of the vertices *huu* and *hdd* listed in the appendix of Ref. [12], and using relation (2.17), we obtain up to $\mathcal{O}\left(v_{SM}^2/f^2\right)$

$$\frac{g_{hff}}{g_{hff}^{SM}} = 1 - \frac{1}{2}\frac{v_{SM}^2}{f^2}\left[\frac{7}{4} + \frac{x^2}{4} - x\right] \qquad f \equiv u, d, c, s, b \tag{2.46}$$

where

$$g_{hff} = \frac{m_f}{v}\, y_f, \qquad g_{hff}^{SM} = \frac{m_f}{v}\bigg|_{v=v_{SM}}. \tag{2.47}$$

Eq. (2.46) will allow us to calculate the tree-level decays of the Higgs boson into two fermions, cf. Appendix A.

Littlest Higgs with T-parity

To implement T-parity in the fermion sector one introduces two $SU(2)_A$ fermion doublets $q_A = (id_{L_A}, -iu_{L_A})^T$ with $A = 1, 2$, as in [20]: T-parity will be defined such that $q_1 \leftrightarrow -q_2$. The T-even combination $u_{L+} = (u_{L_1} - u_{L_2})/\sqrt{2}$ will be the *up*-type component of the SM fermion doublet, while the T-odd combination $u_{L-} = (u_1 + u_2)/\sqrt{2}$ will be its T-odd partner: the same definitions hold also for the *down*-type components.

We require that the T-even (SM) eigenstates obtain a mass only from Yukawa-like interactions after EWSB, while forcing the masses of the T-odd eigenstates to be at the TeV scale. A possible Lagrangian that could generate a TeV mass only for the T-odd combinations can be found in [20]:

$$\mathcal{L}_k \supset -\sqrt{2}kf\left[\bar{d}_{L-}\tilde{d}_c + \frac{1+c_\xi}{2}\bar{u}_{L-}\tilde{u}_c - \frac{s_\xi}{\sqrt{2}}\bar{u}_{L-}\chi_c - \frac{1-c_\xi}{2}\bar{u}_{L-}u_c\right] + $$
$$- m_q\,\bar{u}_c'\,u_c - m_q\,\bar{d}_c'\,d_c - m_\chi\,\bar{\chi}_c'\,\chi_c \quad + \text{h.c.} \tag{2.48}$$

where $c_\xi = \cos\left(h/\sqrt{2}f\right)$, $s_\xi = \sin\left(h/\sqrt{2}f\right)$. u_{L-} and d_{L-} are the T-odd eigenstates, while the other fields $u_c, d_c, \tilde{u}_c, \tilde{d}_c, u_c', d_c', \chi_c, \chi_c'$ are all embedded in the so called *mirror fermions* necessary to write down an invariant Lagrangian under all symmetries. k, m_q and m_χ are

matrices in flavor space for both quarks and leptons: we will assume for simplicity that these matrices are all diagonal and flavor independent.

One can notice that the *down*-type fermions have only Dirac-mass terms and no interactions with the Higgs:

$$- \sqrt{2} k f \, \bar{d}_{L-} \, \tilde{d}_c - m_q \, \bar{d}'_c \, d_c. \tag{2.49}$$

They are thus already mass eigenstates with masses

$$m_1 = \sqrt{2} k f, \qquad m_2 = m_q \tag{2.50}$$

and their contributions will not be considered in the effective one-loop couplings of the Higgs, since they do not couple to the Higgs at tree level.

On the other side, the *up*-type combinations in (2.48) have Dirac-mass terms and also couplings with the Higgs (c_ξ and s_ξ): by diagonalizing these couplings, one obtains the following mass spectrum at leading order

$$m_1^h = \sqrt{2} k f, \qquad m_2^h = m_\chi, \qquad m_3^h = m_q \tag{2.51}$$

where the superscript h indicates that the eigenstates also have an interaction with the Higgs field. The resulting couplings with the Higgs up to $\mathcal{O}\left(v^2/f^2\right)$ are

$$\mathcal{L}_k \supset - \frac{m_1^h}{v} \, y_1 \, \bar{u}_1 \, u_1 \, h - \frac{m_2^h}{v} \, y_2 \, \bar{u}_2 \, u_2 \, h - \frac{m_3^h}{v} \, y_3 \, \bar{u}_3 \, u_3 \, h \tag{2.52}$$

with

$$y_1 = -\frac{1}{4} \frac{v^2}{f^2} \frac{1}{1 - \frac{2 f^2 k^2}{m_\chi^2}}, \qquad y_2 = -\frac{k^2 v^2}{2 m_\chi^2} \frac{1}{1 - \frac{2 f^2 k^2}{m_\chi^2}}, \qquad y_3 = \mathcal{O}\left(\frac{v^4}{f^4}\right). \tag{2.53}$$

We can further reduce the number of free parameters assuming that m_q and m_χ are large enough such that the Higgs couplings (2.53) are independent from their values up to $\mathcal{O}\left(v^2/f^2\right)$, i.e.

mass (*up*-type)	Higgs coupling (only *up*-type)	mass (*down*-type)
$m_{1,i}^h = \sqrt{2} k f$	$y_1^i = -\frac{1}{4} \frac{v^2}{f^2}$	$m_{1,i} = \sqrt{2} k f$
$m_{2,i}^h = m_\chi$	$y_2^i = 0$	$m_{2,i} = m_q$
$m_{3,i}^h = m_q$	$y_3^i = \mathcal{O}\left(v^4/f^4\right)$	

where we have restored the flavor index $i = 1, 2, 3$ referring to both quarks and leptons.

Under these assumptions, in the effective one-loop couplings of the Higgs we will then consider only the contributions from the three degenerate *up*-type T-odd quarks u_1^i, since the other couplings are either suppressed (y_2^i, y_3^i) or absent (*down*-type). The T-odd heavy neutrinos coming from the same interactions are clearly not included in the couplings of the Higgs with gluons and photons, being neither colored, nor electrically charged.

These new twelve T-odd partners u_1^i, d_1^i of the SM fermions can also generate four-fermion operators via box diagrams involving the exchange of NGBs [18]. Assuming always that the couplings k are flavor-diagonal and flavor-independent, the generated operators have the form

$$\mathcal{O}_{\text{4-f}} = -\frac{k^2}{128\,\pi^2 f^2}\bar{\psi}_L\gamma^\mu\psi_L\bar{\psi}'_L\gamma_\mu\psi'_L + \mathcal{O}\left(\frac{g}{k}\right), \tag{2.54}$$

where ψ and ψ' are (distinct) SM fermions. The experimental bound on four-fermion interactions involving SM fields provides an upper bound on the T-odd fermion masses: the strongest constraint comes from the *eedd* operator, whose coefficient is required to be smaller than $2\pi/(26.4\text{ TeV})^2$ [18, 49], which thus yields

$$k^2 \lesssim 0.367\,\pi^3\,f_{\text{TeV}}^2 \tag{2.55}$$

where f_{TeV} is the value of f in units of TeV. Taking a closer look to the contribution of the T-odd fermions u_1^i to the signal strength modifier, one can notice that their contribution enters only in the combination

$$F_{1/2}(m_1^h) \cdot y_1 \tag{2.56}$$

in the expression of the partial decay widths of the Higgs into two gluons and photons, cf. Eq. ($A.6$) and ($A.7$), respectively. However the coupling y_1 is independent of k at the order we are considering, and the loop factor $F_{1/2}$ approaches a constant value $F_{1/2} \to -4/3$ when the particle in the loop is much heavier than the Higgs [25], as in our case ($m_1^h \gg m_h$): the net contribution of the heavy T-odd fermions is thus in good approximation independent of the value of k. Without loss of generality we could therefore choose k to saturate the four-fermion interaction bound (2.55), with an upper limit of 4π when $f \to \infty$.

The next task is to write invariant Yukawa-like terms to give mass to the T-even (SM) combinations u_{L+} and d_{L+}. In order to avoid dangerous contributions to the Higgs mass from one-loop quadratic divergences, the top Yukawa sector must also incorporate a collective symmetry breaking pattern.

The details of the procedure could again be found in [20]: the Yukawa-like Lagrangian for the top sector contains the following terms

$$\mathcal{L}_t \supset -\lambda_1 f\left(\frac{s_\Sigma}{\sqrt{2}}\,\bar{t}_{L+}\,t'_R + \frac{1+c_\Sigma}{2}\,\bar{T}'_{L+}\,t'_R\right) - \lambda_2 f\left(\bar{T}'_{L+}\,T'_{R+} + \bar{T}'_{L-}\,T'_{R-}\right) + \text{h.c.} \tag{2.57}$$

where $c_\Sigma = \cos\left(\sqrt{2}h/f\right)$ and $s_\Sigma = \sin\left(\sqrt{2}h/f\right)$.

Among the terms that we have neglected, there are the interaction terms of the T-odd eigenstate t_{L-}, which does not acquire any mass term from $\mathcal{L}_t$ while obtaining its mass from $\mathcal{L}_k$ as explained before. In $\mathcal{L}_t$ a different T-odd Dirac fermion $T_- \equiv \left(T'_{L-}, T'_{R-}\right)$ obtains a high-scale mass

$$m_{T_-} = \lambda_2\, f\ . \tag{2.58}$$

It does not have tree-level interactions to the Higgs boson, and will be thus not included in the Higgs one-loop effective couplings.

One should notice that the T-even top Lagrangian (2.57) has the same form as the L^2H top Lagrangian (2.41): the mass spectrum and the couplings to the Higgs boson will

therefore be the same in both models, by simply setting $x = 0$. The T-even combinations in $\mathcal{L}_t$, i.e. (t_{L+}, t'_R) and (T'_{L+}, T'_{R+}), mix among each other:

$$-\mathcal{L}_t \supset \begin{pmatrix} \bar{t}_{L+} & \bar{T}'_{L+} \end{pmatrix} \mathcal{M} \begin{pmatrix} t'_R \\ T'_{R+} \end{pmatrix} + \text{h.c.}, \qquad \mathcal{M} = \begin{pmatrix} \frac{\lambda_1 f}{\sqrt{2}} \sin \frac{\sqrt{2}h}{f} & 0 \\ \lambda_1 f \cos^2 \frac{h}{\sqrt{2}f} & \lambda_2 f \end{pmatrix}. \tag{2.59}$$

The mass terms are diagonalized by defining the linear combination [18]

$$t_L = \cos\beta \cdot t_{L+} - \sin\beta \cdot T'_{L+}, \qquad T_{L+} = \sin\beta \cdot t_{L+} + \cos\beta \cdot T'_{L+}$$
$$t_R = \cos\alpha \cdot t'_R - \sin\alpha \cdot T'_{L+} \qquad T_{R+} = \sin\alpha \cdot t'_R + \cos\alpha \cdot T'_{R+} \tag{2.60}$$

Here, we use the dimensionless ratio $R = \lambda_1/\lambda_2$ as well as the leading order expressions of the mixing angles

$$\sin\alpha = \frac{R}{\sqrt{1 + R^2}}, \qquad \sin\beta = \frac{R^2}{1 + R^2} \frac{v}{f}. \tag{2.61}$$

The leading order mass spectrum is the following

$$m_t = \frac{\lambda_2 R}{\sqrt{1 + R^2}} v, \qquad m_{T_+} = \lambda_2 \sqrt{1 + R^2} f \quad . \tag{2.62}$$

Again, R and λ_2 are considered to be free parameters. However we can fix λ_2 requiring that, for given (f, R), m_t corresponds to the experimental top mass value: this way, the only free parameters in the T-even top sector are f and R.

The resulting couplings to the Higgs up to $\mathcal{O}\left(v^2/f^2\right)$ are given by [14, 16]

$$\mathcal{L}_t \supset -\frac{m_t}{v} y_t \, \bar{t}\, t\, h - \frac{m_{T_+}}{v} y_{T_+} \, \bar{T}_+ \, T_+ \, h \tag{2.63}$$

with

$$y_t = 1 - \frac{v^2}{f^2} \frac{2R^4 + R^2 + 2}{3\left(1 + R^2\right)^2}, \qquad y_{T_+} = -\frac{v^2}{f^2} \frac{R^2}{\left(1 + R^2\right)^2}. \tag{2.64}$$

The other two generations of T-even (SM-like) *up*-type quarks acquire their mass through analogous terms as $\mathcal{L}_t$, but with the $T_\pm$ missing since the Yukawa couplings are small and one does not have to worry about the quadratic divergences. Using Eq. (2.20), the corrections to the Yukawa couplings with respect to their SM values up to $\mathcal{O}\left(v_{SM}^4/f^4\right)$ are given by [20]

$$\frac{g_{h\bar{u}u}}{g_{h\bar{u}u}^{SM}} = 1 - \frac{3}{4} \frac{v_{SM}^2}{f^2} - \frac{5}{32} \frac{v_{SM}^4}{f^4} \qquad u \equiv u, c. \tag{2.65}$$

We need to construct also a Yukawa interaction which gives a mass after EWSB to the T-even (SM-like) *down*-type quarks and charged leptons. Two possible constructions of Lagrangians can be found in [20], which will be denoted as *Case A* and *Case B*, respectively. The corresponding corrections to the Yukawa couplings with respect to their SM values up to $\mathcal{O}\left(v_{SM}^4/f^4\right)$ are given by ($d \equiv d, s, b, l_i^\pm$)

$$\frac{g_{h\bar{d}d}}{g_{h\bar{d}d}^{SM}} = 1 - \frac{1}{4} \frac{v_{SM}^2}{f^2} + \frac{7}{32} \frac{v_{SM}^4}{f^4} \qquad \text{Case A}$$

$$\frac{g_{h\bar{d}d}}{g_{h\bar{d}d}^{SM}} = 1 - \frac{5}{4} \frac{v_{SM}^2}{f^2} - \frac{17}{32} \frac{v_{SM}^4}{f^4} \qquad \text{Case B.} \tag{2.66}$$

We will analyze the parameter space of the *LHT* model with both *Case A,B* implementations: it is to be noted that *Case B* predicts a stronger suppression for the down-type fermion couplings to the Higgs boson, and this will have an influence on our results.

Simplest Little Higgs

Since this model contains a gauged $SU(3)$, SM fermions that are doublets under $SU(2)$ must be enlarged into triplets under $SU(3)$. In addition, new $SU(3)$ singlet fermions must be introduced to cancel the hypercharge anomalies and to give mass to the new third components of the $SU(3)$ triplet fermions. In the "anomaly-free" scenario, the quarks of the third generation and all leptons are embedded into **3** of $SU(3)$:

$$Q_3^T = (t, b, iT), \qquad L_m^T = (\nu_m, l_m, iN_m) \quad (m = 1, 2, 3) \tag{2.67}$$

adding also the corresponding right handed singlets it^c, ib^c, iT^c and il_m^c, iN_m^c. We do not include a right-handed neutrino, leaving the neutrinos as massless.

The corresponding Yukawa Lagrangian $\mathcal{L}_Y$ can be found explicitly in [21, 23], and gives rise at leading order to the following mass spectrum after EWSB: The free parameters

$$m_b \propto \lambda^b \qquad m_t = \lambda_2^t\, v\, R\sqrt{\frac{t_\beta^2+1}{2(t_\beta^2+R^2)}} \qquad m_T = \lambda_2^t\, f\sqrt{\frac{t_\beta^2+R^2}{t_\beta^2+1}}$$

$$m_\nu = 0 \qquad\qquad m_{l_m} \propto \lambda_{m,n}^l \qquad\qquad m_{N_m} = \lambda_{N_m}\, f\, \frac{t_\beta}{\sqrt{1+t_\beta^2}}$$

are $R = \left(\lambda_1^t/\lambda_2^t\right)$, λ_2^t, λ^b f, t_β in the top sector, and λ^{N_m}, $\lambda_{m,n}^l$, f, t_β in the lepton sector, respectively. However we can fix λ_2^t, λ^b and $\lambda_{m,n}^l$ requiring that for given (f, t_β) the predicted values of m_t, m_b and m_{l_m} correspond to their experimental values: in this way the free parameters in the top sector are f and t_β (notice that R is fixed by the EWSB requirement), while in the lepton sector they are λ_{N_m}, f, t_β. Since the heavy neutrinos do not affect neither the effective *hgg* coupling, nor the *hγγ* one, the parameter λ_{N_m} is thus irrelevant for our study.

Regarding the couplings of these fermions to the Higgs, they can be parametrized as

$$\mathcal{L}_Y \supset -\frac{m_l^i}{v}\, y_l^i\, \bar{l}_i\, l_i\, h - \frac{m_N^i}{v}\, y_N^i\, \bar{N}_i\, N_i\, h - \frac{m_t}{v}\, y_t\, \bar{t}\, t\, h - \frac{m_T}{v}\, y_T\, \bar{T}\, T\, h - \frac{m_b}{v}\, y_b\, \bar{b}\, b\, h \tag{2.68}$$

and predicted up to $\mathcal{O}\left(v^2/f^2\right)$ to be [21]

$$y_l^i = 1 - \frac{1}{6}\frac{v^2}{f^2}\left(3 + \frac{t_\beta^4 - t_\beta^2 + 1}{t_\beta^2}\right), \qquad y_N^i = \mathcal{O}\left(\frac{v^4}{f^4}\right)$$

$$y_t = 1 - \frac{1}{6}\frac{v^2}{f^2}\left[\frac{\left(1+t_\beta^2\right)^2\left(R^4 - R^2\, t_\beta^2 + t_\beta^4\right)}{t_\beta^2\left(t_\beta^2 + R^2\right)^2}\right], \qquad y_T = -\frac{1}{2}\frac{v^2}{f^2}R^2\left(\frac{t_\beta^2+1}{t_\beta^2 + R^2}\right)^2$$

$$y_b = 1 - \frac{1}{6}\frac{v^2}{f^2}\left(3 + \frac{t_\beta^4 - t_\beta^2 + 1}{t_\beta^2}\right). \tag{2.69}$$

The corrections of the bottom-quark and lepton Yukawa couplings with respect to their SM values up to $\mathcal{O}\left(v_{SM}^4/f^4\right)$ are equivalently given by

$$\frac{g_{h\bar{b}b}}{g_{h\bar{b}b}^{SM}} = \frac{g_{h\bar{l}l}}{g_{h\bar{l}l}^{SM}} = 1 - \frac{1}{4}\frac{v_{SM}^2}{f^2}\left(\frac{t_\beta^4 + t_\beta^2 + 1}{t_\beta^2}\right) - \frac{1}{720}\frac{v_{SM}^4}{f^4}\left(\frac{t_\beta^8 + 24t_\beta^6 - 19t_\beta^4 + 24t_\beta^2 + 1}{t_\beta^4}\right) \tag{2.70}$$

where as usual

$$g_{hff} = \frac{m_f}{v}\, y_f, \qquad g_{hff}^{SM} = \frac{m_f}{v}\Big|_{v=v_{SM}}. \tag{2.71}$$

In the "anomaly free" embedding, the first two generations of quarks are embedded into $\mathbf{3}^*$ of $SU(3)$ with the corresponding right-handed singlets:

$$Q_1^T = (d, -u, iD) \qquad id^c, iu^c, iD^c$$
$$Q_2^T = (s, -c, iS) \qquad is^c, ic^c, iS^c$$

Notice that the heavy vector-like quarks of the first two generations have electric charge $-1/3$ in contrast to the charge $+2/3$ of the heavy quark of the third generation. The Lagrangian terms for the Yukawa couplings of the first two generations of quarks can be found in [21], and the resulting mass spectrum after EWSB consists of the SM quarks plus two heavy D, S quarks with charge $-1/3$.

As suggested in [21], one would expect an $h\, D_m^c\, D_m$ coupling at order v/f as for the top sector, but this term is exactly canceled by the contribution from $h\, D_m^c\, d_m$ after d-D mixing if the down and strange quark masses are neglected. For this reason we will not include the heavy D, S in the calculation of the one-loop effective couplings: therefore only the contributions from the top (the dominant one among the SM-like particles) and from the heavy top T will be included in the one-loop effective couplings.

The corrections of the charm and strange quarks Yukawa couplings with respect to their SM values up to $\mathcal{O}\left(v_{SM}^4/f^4\right)$ are finally given by

$$\frac{g_{h\bar{c}c}}{g_{h\bar{c}c}^{SM}} = \frac{g_{h\bar{s}s}}{g_{h\bar{s}s}^{SM}} = 1 - \frac{1}{4}\frac{v_{SM}^2}{f^2}\left(\frac{t_\beta^4 + t_\beta^2 + 1}{t_\beta^2}\right) - \frac{1}{720}\frac{v_{SM}^4}{f^4}\left(\frac{t_\beta^8 + 24t_\beta^6 - 19t_\beta^4 + 24t_\beta^2 + 1}{t_\beta^4}\right). \tag{2.72}$$

2.3 Electroweak Precision Observables

Littlest Higgs

The contribution to EWPO from the L^2H structure was already studied in Ref. [15]. They calculated all contributions to EWPO from the tree-level exchange of the heavy gauge bosons and from the presence of the triplet *vev* v': there should in principle be also contributions due to heavy quark loop modifications to the light gauge boson propagators, but as the authors of Ref. [15] have shown, these contributions are almost an order of magnitude smaller than the maximal contribution of the triplet *vev*. Therefore we ignore them.

We refer to the original Ref. [15] for the explicit expression of the 21 EWPO in terms of the parameters f, c, c', x defined above (the list of the 21 EWPO can be found in Sec. 3). We adapt their notation as follows:

$$\Delta = \frac{v^2}{f^2}, \qquad \Delta' = \frac{x^2}{16}\frac{v^2}{f^2}. \tag{2.73}$$

A derivation of the oblique parameters in the L^2H model can be found in Ref. [26].

Littlest Higgs with T-parity

Due to the introduction of T-parity, no T-odd state can contribute as external state at tree-level: therefore no contributions to electroweak observables arise at tree-level from T-odd states. The only new particle which is T-even is the T-even top partner T_+, but it can contribute at tree level only to observables involving the SM top quark, such as its couplings to W and Z bosons: since these couplings have not been measured experimentally yet, no constraints arise at tree-level also from the T-even top partner. We will then consider only the one-loop contributions to EWPO coming from the new T-even/odd states, using the results of Ref. [18, 28, 29].

At one loop, oblique corrections to the electroweak gauge boson propagators induced by diagrams involving the top and its T_+ partner are given by

$$
\begin{aligned}
S_{T_+} &= \frac{s_\beta^2}{2\pi}\left[\left(\frac{1}{3}-c_\beta^2\right)\log x_t + c_\beta^2\frac{(1+x_t)^2}{(1-x_t)^2} + \frac{2c_\beta^2 x_t^2(3-x_t)\log x_t}{(1-x_t)^3} - \frac{8c_\beta^2}{3}\right] \\
T_{T_+} &= \frac{3}{16\pi}\frac{s_\beta^2}{s_w^2 c_w^2}\frac{m_t^2}{m_Z^2}\left[\frac{s_\beta^2}{x_t} - 1 - c_\beta^2 - \frac{2c_\beta^2}{1-x_t}\log x_t\right] \\
U_{T_+} &= -\frac{s_\beta^2}{2\pi}\left[s_\beta^2\log x_t + c_\beta^2\frac{(1+x_t)^2}{(1-x_t)^2} + \frac{2c_\beta^2 x_t^2(3-x_t)\log x_t}{(1-x_t)^3} - \frac{8c_\beta^2}{3}\right]
\end{aligned} \tag{2.74}
$$

Here, $s_\beta = \sin\beta$ is the mixing angle in the right-handed top sector, $x_t = m_t^2/m_{T_+}^2$, and s_w is the sine of the Weinberg angle.

The T-odd top partner T_- does not contribute to the S,T,U parameters since it is an $SU(2)_L$ singlet which does not mix with the SM top. Moreover, the corrections from T_- loops are very small, and we do not include them in our fit. But the other T-odd heavy fermions coming from the interactions in Eq. (2.48) give a contribution to the T parameter at one-loop: under the assumption of flavor-independent k, the contribution of each T-odd fermion partner is given up to $\mathcal{O}\left(v^2/f^2\right)$ by

$$T_{\text{T-odd}} = -\frac{k^2}{192\pi^2\alpha_w}\frac{v^2}{f^2}. \tag{2.75}$$

As explained in detail in the previous section, there is an upper bound on the value of k coming from four-fermion interactions involving SM fields, Eq. (2.55): the maximum contribution to the T parameter consistent with this bound becomes

$$|T_{\text{T-odd}}| \lesssim 0.05 \tag{2.76}$$

for each T-odd fermion partner. The large number of T-odd partners (twelve) could thus have a sizable effect on the EWPO: however a smaller value of k could reduce this contribution, and so we have not included it in our fit.

Regarding the contribution from the gauge sector, the authors of Ref. [29] calculated that the total log-divergent contribution due to the custodial $SU(2)$-violating tree-level mass splitting of the T-odd heavy W_H^3 and $W_H^\pm$ gauge bosons completely vanishes, leaving only negligible finite terms of order v/f which we do not include in our analysis.

Another important correction to both the S and T parameters follows from the modified couplings of the Higgs boson to the SM gauge bosons. In the SM, due to its renormalizability, the one-loop contribution of the Higgs boson to the vector boson self energies exactly cancels the logarithmic divergence arising from loops of would-be NGBs in the gauge-less limit [30]. As first noticed by the authors of Ref. [31], the modified Higgs couplings to the SM gauge bosons imply that the contribution of the Higgs to the self-energy does not exactly cancel the infrared log-divergence arising from the NGBs, leading to the following contributions to the oblique parameters:

$$
\begin{aligned}
S_h &= -\frac{1}{6\pi}\left(1 - y_W^2\right)\log\frac{m_h}{\Lambda} \\
T_h &= \frac{3}{8\pi c_w^2}\left(1 - y_W^2\right)\log\frac{m_h}{\Lambda} \quad .
\end{aligned}
\tag{2.77}
$$

y_W parametrizes the shift of the coupling of one Higgs boson to the SM gauge bosons in the usual notation, and $\Lambda = 4\pi f$ is the cut-off of the non-linear sigma model. In particular for the LHT model we obtain:

$$
\begin{aligned}
S_h &= -\frac{1}{18\pi}\frac{v^2}{f^2}\log\frac{m_h}{\Lambda} \\
T_h &= \frac{1}{8\pi c_w^2}\frac{v^2}{f^2}\log\frac{m_h}{\Lambda}.
\end{aligned}
\tag{2.78}
$$

The LHT model contains an additional T-odd $SU(2)_L$-triplet scalar field ϕ: the effects on the S, T, U parameters are of order $\mathcal{O}(v^4/m_\phi^4)$ and therefore negligible for m_ϕ in its natural range, around 1 TeV. We will thus not include these effects in our fit.

Finally, other possible contributions arise from new operators which parametrize the effects of the UV physics on weak-scale observables:

$$
\begin{aligned}
S_{UV} &= c_s\frac{4m_W^2}{\pi g^2 f^2} \\
T_{UV} &= -c_t\frac{m_W^2}{2\pi e^2 g^2 f^2} \quad ,
\end{aligned}
\tag{2.79}
$$

where c_s and c_t are again coefficients of order one whose exact values depend on the details of the UV physics, and which for simplicity we assume to be equal to one as in [28]. All these different contributions to the oblique parameters are then summed up.

The only important non-oblique correction to the neutral-current interactions which could affect the EWPO is the one-loop T_+ contribution to the $Zb_L\bar{b}_L$ vertex: to leading

order in the limit $m_{T+} \gg m_t \gg m_W$ it is given by

$$\delta \tilde{g}_L^{b\bar{b}} = \frac{g}{c_w} \frac{\alpha_w}{8\pi s_w^2} \frac{m_t^4}{m_W^2 \, m_{T+}^2} R^2 \log \frac{m_{T+}^2}{m_t^2}. \tag{2.80}$$

where we have used the notation of Appendix B where more details on the calculation of the EWPO can be found.

With the explicit expressions of the oblique parameters S, T, U and of the neutral-current coefficient $\delta \tilde{g}_L^{b\bar{b}}$, we can finally obtain the explicit expressions of the 21 EWPO using the general results of Ref. [32] summarized in Appendix B. The list of these variables and their experimental values can be found in Sec. 3.

Simplest Little Higgs

The dominant tree-level contributions to the oblique parameters in the *SLH* model come from the presence of a Z' boson with $Z - Z'$ mixing: their explicit expression can be found in Ref. [8, 33] as

$$S_{Z'} = \frac{8 s_w^2}{\alpha_w} \frac{m_W^2}{g^2 f^2}, \qquad T_{Z'} = \frac{1}{\alpha_w} \cdot \frac{v^2}{8 f^2} (1 - t_w^2)^2. \tag{2.81}$$

The oblique parameters receive contributions also from the modification in the Higgs couplings to the electroweak gauge bosons w.r.t. their SM values, cf. Eq. (2.77):

$$\begin{aligned} S_h &= -\frac{1}{18\pi} \frac{t_\beta^4 - t_\beta^2 + 1}{t_\beta^2} \frac{v^2}{f^2} \log \frac{m_h}{\Lambda} \\ T_h &= \frac{1}{8\pi c_w^2} \frac{t_\beta^4 - t_\beta^2 + 1}{t_\beta^2} \frac{v^2}{f^2} \log \frac{m_h}{\Lambda} \end{aligned} \qquad . \tag{2.82}$$

Again, $\Lambda = 4\pi f$ is the cut-off of the non-linear sigma model.

The corrections in the flavor sector can be read off from the fermion-gauge interaction Lagrangians in [21]: following their assumptions, we will ignore right-handed mixing and choose the Yukawa parameters in order to suppress the heavy-light mixing effects in the first and second generations of quarks and in the b-quark sector[1]. Using the notation of Appendix B and defining the quantities

$$\delta_\nu = -\frac{v}{\sqrt{2} f t_\beta}, \qquad \delta_Z = -\frac{(1 - t_w^2)\sqrt{3 - t_w^2}}{8 c_w} \frac{v^2}{f^2} \quad , \tag{2.83}$$

which parametrize the rotation to the mass eigenstates in the fermion- and in the neutral gauge boson sectors, respectively, the corrections to the charged-current couplings up to $\mathcal{O}\left(v^2/f^2\right)$ are given in Table 1, while the corrections to the neutral-current couplings (with $u \equiv u, c$ and $d \equiv d, s$) are given in Table 2.

With the explicit expressions of the oblique parameters S, T, U and of the charged- and neutral-current coefficients, we can finally obtain the explicit expressions of the 21 EWPO using the general results of Ref. [32] cited in Appendix B.

---

[1]Using the notation of Ref. [21], under these assumptions, we obtain the following relations: $\Delta u_i \simeq \Delta u_3 = V_{33}^{u*} \delta_\nu t_\beta^2 \frac{1 - R^2}{R^2 + t_\beta^2} \simeq \delta_\nu t_\beta^2 \frac{1 - R^2}{R^2 + t_\beta^2}$, $\Delta D d_j \simeq \Delta D d \simeq \delta_\nu$ and $\Delta S d_j \simeq \Delta S s \simeq \delta_\nu$.

	$\delta\tilde{h}_L$	$\delta\tilde{h}_R$
νl	$-\delta_\nu^2/2$	0
ud	$-\delta_\nu^2/2$	0
cs	$-\delta_\nu^2/2$	0
tb	$-\frac{1}{2}\delta_\nu^2 t_\beta^4 \frac{(1-R^2)^2}{\left(R^2+t_\beta^2\right)^2}$	0

Table 1. Corrections of the charged-current couplings in *SLH*.

	$\delta\tilde{g}_L$	$\delta\tilde{g}_R$
$\nu\nu$	$-\delta_\nu^2/2 + (1/2-s_w^2)\delta_Z/\sqrt{3-4s_w^2}$	0
ll	$(1/2-s_w^2)\delta_Z/\sqrt{3-4s_w^2}$	$s_w^2\delta_Z/\sqrt{3-4s_w^2}$
uu	$(-1/2+2/3s_w^2)\delta_Z/\sqrt{3-4s_w^2}$	$-2/3s_w^2\delta_Z/\sqrt{3-4s_w^2}$
tt	$-\frac{1}{2}\delta_\nu^2 t_\beta^4 \frac{(1-R^2)^2}{\left(R^2+t_\beta^2\right)^2} + (1/2-1/3s_w^2)\delta_Z/\sqrt{3-4s_w^2}$	$-2/3s_w^2\delta_Z/\sqrt{3-4s_w^2}$
dd	$\delta_\nu^2/2 + (-1/2+2/3s_w^2)\delta_Z/\sqrt{3-4s_w^2}$	$1/3s_w^2\delta_Z/\sqrt{3-4s_w^2}$
bb	$(1/2-1/3s_w^2)\delta_Z/\sqrt{3-4s_w^2}$	$1/3s_w^2\delta_Z/\sqrt{3-4s_w^2}$

Table 2. Corrections of the neutral-current couplings in *SLH*.

3 Statistical Method and Experimental Data

It is customary for the experimental collaborations to express the results of the SM-like Higgs searches in terms of a *signal strength modifier* μ, defined as the factor by which the SM Higgs signal is modified for a given value of m_h:

$$\mu^i = \frac{n_S^i}{n_S^{SM,i}} = \frac{\sum_p \sigma_p \cdot \epsilon_i^p}{\sum_p \sigma_p^{SM} \cdot \epsilon_i^p} \cdot \frac{BR_i}{BR_i^{SM}} \tag{3.1}$$

where i, p are indices for a specific decay channel and production mode, respectively. ϵ_i^p is the efficiency of the kinematic cuts for a given production mode p and decay channel i, and n_S^i is the number of expected Higgs signal events evaluated in a chosen model (e.g. $n_S^{SM,i}$ is evaluated in the SM).

For each Higgs decay channel considered, the ATLAS and CMS collaborations usually report the 95% *CL* limit on μ ($\mu_{95\%}$) and the best-fit value $\hat{\mu}$ for a given hypothesis on m_h. In particular, values $\mu_{95\%} < 1$ exclude at 95% *CL* the SM Higgs for that particular value of the Higgs mass. The efficiencies of the kinematic cuts are instead not reported by the collaborations (the only exceptions are all the $\gamma\gamma$ and CMS $\tau\tau$ 8 TeV channels), making it thus very hard (if not impossible) to correctly compare a theory prediction with the observed data.

To implement a χ^2 analysis we follow the procedure described in [34]. One defines the covariance matrix $\mathcal{C}$ of the observables, and $\Delta\theta_i$ as the vector of the difference in the observed and predicted value of the observables, which is a function of the free parameters of the model. The χ^2 measure is then given by

$$\chi^2 = (\Delta\theta_i)^T \left(\mathcal{C}^{-1}\right)_{ij} (\Delta\theta_j). \tag{3.2}$$

The 95% and 99% best-fit *CL* regions are then defined by the cumulative distribution function for an appropriate number of degrees of freedom (*d.o.f.*).

3.1 Higgs searches

First, we consider as observables for the χ^2 measure (3.2) the different best-fit values of the signal strength modifiers of all available public data reported by the ATLAS and CMS collaborations for the 7 and 8 TeV Higgs searches. In particular, we include in our analysis the 7 TeV ~ 5 fb^{-1} and 8 TeV ~ 6 fb^{-1} data from the July 2012 publications of both collaborations, and also the latest December 2012 update of up to ~ 13 fb^{-1} of many of the 8 TeV samples.

For our analysis, we have taken the matrix $\mathcal{C}$ to be diagonal with the sum of the square of the 1σ theory and experimental errors as diagonal entries: off-diagonal correlation coefficients are indeed neglected, as correlation coefficients are currently not supplied by the experimental collaborations. As already discussed in [37], the absence of information regarding correlations in fits of Higgs couplings is not a significant limitation, given the current level of statistical uncertainty. The authors in [37] claim indeed that the error on the best fit point assuming zero correlation is less than 1%, at least in the CMS $\gamma\gamma$ final state.

For the experimental errors we use the quoted 1σ errors on the reported signal strength ($\delta\mu_{i,\mathrm{exp}}$), while for theoretical uncertainties we propagate the cross section error $\delta\sigma_i$ as an uncertainty on the signal strength modifier:

$$\delta\mu_{i,\mathrm{th}} = \mu_i \left(\frac{\sqrt{\sum_j r_j^2 \cdot \delta\sigma_j^2}}{\sum r_j \cdot \sigma_j} - \frac{\sqrt{\sum_j \delta\sigma_j^2}}{\sum_j \sigma_j} \right) \quad , \tag{3.3}$$

where r_j is the appropriate rescaling factor for each cross section j. In this way, the χ^2 measure reduces to the usual form

$$\chi^2 = \sum_i \frac{(\mu_i - \hat{\mu}_i)^2}{\sigma_i^2} \tag{3.4}$$

where μ_i is the i-th signal strength predicted by the model as a function of the free parameters, $\hat{\mu}_i$ is the respective best-fit value, and $\sigma_i = \sqrt{\delta\mu_{i,\mathrm{exp}}^2 + \delta\mu_{i,\mathrm{th}}^2}$ the total uncertainty.

We summarize in Table 3 the available data, reporting in particular the different best-fit values $\hat{\mu}_i$ and the reference masses at which the single $\hat{\mu}_i$ are evaluated. Notice that these reference masses are not necessarily the masses at which the highest local significance has been obtained for each channel. We have chosen the best-fit values of the signal strengths and the corresponding masses in order to be able to reconstruct separately (for which only 7 TeV and combined results are given by the experiments) the 7 and 8 TeV signal strengths. For example, even if the highest significance in the 7+8 TeV combined analysis of the ATLAS ZZ channel [41] has been found for a Higgs mass of 123.5 GeV ($\hat{\mu} = 1.3 \pm 0.5$), we used the 7+8 TeV signal strength for a Higgs mass of 126 GeV ($\hat{\mu} = 0.8 \pm 0.4$) for reconstructing the 8 TeV signal strength reported in Table 3, in order to comply with the 7 TeV signal strength which has been given in Ref. [2] for $m_h = 126$ GeV.

ATLAS 7 TeV	m_h [GeV]	$\hat{\mu}$
$\gamma\gamma_{\text{UClPT}}$ [38]	126.5	0.5 ± 1.5
$\gamma\gamma_{\text{UChPT}}$ [38]	126.5	0.2 ± 2.0
$\gamma\gamma_{\text{URlPT}}$ [38]	126.5	2.5 ± 1.7
$\gamma\gamma_{\text{URhPT}}$ [38]	126.5	10.4 ± 3.7
$\gamma\gamma_{\text{CClPT}}$ [38]	126.5	6.1 ± 2.6
$\gamma\gamma_{\text{CChPT}}$ [38]	126.5	-4.4 ± 1.8
$\gamma\gamma_{\text{CRlPT}}$ [38]	126.5	2.7 ± 2.0
$\gamma\gamma_{\text{CRhPT}}$ [38]	126.5	-1.6 ± 2.9
$\gamma\gamma_{\text{CT}}$ [38]	126.5	0.3 ± 3.6
$\gamma\gamma_{\text{jj}}$ [38]	126.5	2.7 ± 1.9
ZZ [2]	126.0	1.4 ± 1.1
WW [2]	126.0	0.5 ± 0.6
bb [40]	125.0	-2.7 ± 1.6
$\tau\tau$ [2]	126.0	0.4 ± 1.8

ATLAS 8 TeV	m_h [GeV]	$\hat{\mu}$
$\gamma\gamma_{\text{UClPT}}$ [39]	126.5	1.0 ± 0.9
$\gamma\gamma_{\text{UChPT}}$ [39]	126.5	0.3 ± 1.2
$\gamma\gamma_{\text{URlPT}}$ [39]	126.5	2.9 ± 1.2
$\gamma\gamma_{\text{URhPT}}$ [39]	126.5	1.8 ± 1.4
$\gamma\gamma_{\text{CClPT}}$ [39]	126.5	1.5 ± 1.2
$\gamma\gamma_{\text{CChPT}}$ [39]	126.5	1.0 ± 1.6
$\gamma\gamma_{\text{CRlPT}}$ [39]	126.5	2.3 ± 1.2
$\gamma\gamma_{\text{CRhPT}}$ [39]	126.5	0.5 ± 1.6
$\gamma\gamma_{\text{CT}}$ [39]	126.5	2.0 ± 2.0
$\gamma\gamma_{\text{2jhm}}$ [39]	126.5	2.0 ± 1.1
$\gamma\gamma_{\text{2jlm}}$ [39]	126.5	3.6 ± 2.1
$\gamma\gamma_{\text{LT}}$ [39]	126.5	1.2 ± 2.2
ZZ [41]*	126.0	0.7 ± 0.4
WW [42]	126.0	1.4 ± 0.6
bb [40]	125.0	1.0 ± 1.4
$\tau\tau$ [42]*	126.0	0.7 ± 0.8

CMS 7 TeV	m_h [GeV]	$\hat{\mu}$
$\gamma\gamma_{\text{cat0}}$ [43]	125.0	3.2 ± 1.8
$\gamma\gamma_{\text{cat1}}$ [43]	125.0	0.7 ± 0.9
$\gamma\gamma_{\text{cat2}}$ [43]	125.0	0.7 ± 1.2
$\gamma\gamma_{\text{cat3}}$ [43]	125.0	1.5 ± 1.6
$\gamma\gamma_{\text{jj}}$ [43]	125.0	4.2 ± 2.1
ZZ [44]	125.0	0.6 ± 0.6
WW [44]	125.0	0.4 ± 0.6
bb [44]	125.0	0.6 ± 1.2
$\tau\tau$ [48]	125.0	1.0 ± 0.9

CMS 8 TeV	m_h [GeV]	$\hat{\mu}$
$\gamma\gamma_{\text{cat0}}$ [43]	125.0	1.4 ± 1.2
$\gamma\gamma_{\text{cat1}}$ [43]	125.0	1.5 ± 1.0
$\gamma\gamma_{\text{cat2}}$ [43]	125.0	0.9 ± 1.2
$\gamma\gamma_{\text{cat3}}$ [43]	125.0	3.8 ± 1.8
$\gamma\gamma_{\text{jj loose}}$ [43]	125.0	-0.6 ± 2.0
$\gamma\gamma_{\text{jj tight}}$ [43]	125.0	1.3 ± 1.6
ZZ [45]*	125.0	0.9 ± 0.4
WW [46]*	125.0	0.8 ± 0.3
bb [47]*	125.0	1.5 ± 0.7
$\tau\tau$ [48]	125.0	0.6 ± 0.6

Table 3. Signal-strength best-fit values of the 7 and 8 TeV samples collected by ATLAS and CMS. The red asterisks mark the 8 TeV channels for which the best-fit signal-strengths have been reconstructed from the 7 and 7+8 TeV values.

For some channels, indeed only the combined 7+8 TeV signal strengths are available in addition to the 7 TeV results. As suggested in [34], one can assume a Gaussian approximation for the probability density functions (*pdf*) describing the different signal strengths, i.e.

$$p[\mu_i|\hat{\mu}_i, \sigma_i] \simeq e^{(\mu_i - \hat{\mu}_i)^2/(2\sigma_i^2)} \tag{3.5}$$

and obtain the combined *pdf* by multiplying the individual channel *pdf*s. The combined

pdf is therefore also Gaussian, with central value $\hat{\mu}_c$ and width σ_c approximately given by

$$\frac{1}{\sigma_c^2} = \sum_i \frac{1}{\sigma_i^2}, \qquad \frac{\hat{\mu}_c}{\sigma_c^2} = \sum_i \frac{\hat{\mu}_i}{\sigma_i^2}. \tag{3.6}$$

By solving the previous equations, one can then reconstruct the unknown 8 TeV data from the reported 7 and 7+8 TeV data, as has been done for the channels marked with a red asterisk in Table 3.

It is to be noted that the collaborations have reported the best-fit values of the 7 and 8 TeV diphoton channels exclusively with respect to the different selection cut categories, and reported also the different cut efficiencies for the single subchannels, cf. Ref. [38, 43]. Therefore we were able to add all single diphoton subchannel contributions to the χ^2 measure, exclusively with respect to the production modes.

For the $h \rightarrow bb$ channels we assume the results to be fully dominated by the Higgs-Strahlung production mode, neglecting the contributions from the other production modes. For all other channels we considered the signal as inclusive with respect to the production modes, neglecting thus the cut efficiencies, since they have not been reported.

The SM parameters have been obtained from the updated values of the Particle Data Group Collaboration [49], while for the SM Higgs production cross sections with respective uncertainties and branching ratios we have used the recommended values by the *LHC Higgs Cross Section Working Group* [50].

3.2 Electroweak Precision Data

We incorporate EWPO by directly adding their contribution to the χ^2 measure. In particular, we include the contribution from the following 21 different low-energy and Z-pole precision observables for $m_h = 124.5$ GeV [49], as summarized in Table 4.

Since no correlation coefficients are supplied for these 21 observables, we will assume them as independent and add their contribution to the χ^2 measure as

$$\chi^2 = \sum_i \frac{(\mathcal{O}_i - \hat{\mathcal{O}}_i)^2}{\sigma_i^2} \qquad . \tag{3.7}$$

Here, $\mathcal{O}_i$ is the i-th observable predicted by the model as a function of the free parameters, $\hat{\mathcal{O}}_i$ is the respective measured value, and σ_i the experimental uncertainty.

4 Results

In order to calculate the updated exclusion contours for the different models, we need to determine the explicit expression of the signal strength modifier μ^i

$$\mu^i = \frac{\sum_p \sigma_p^{\text{LH}} \cdot \epsilon_i^p}{\sum_p \sigma_p^{\text{SM}} \cdot \epsilon_i^p} \cdot \frac{BR_i^{\text{LH}}}{BR_i^{SM}} \tag{4.1}$$

for each decay channel i. μ^i depends on the different free parameters of the model under which it is evaluated: in particular, the three models we are considering share a free

	Value	SM prediction
Γ_Z [GeV]	2.4952 ± 0.0023	2.4961 ± 0.0010
R_e	20.804 ± 0.050	20.744 ± 0.011
R_μ	20.785 ± 0.033	20.744 ± 0.011
R_τ	20.764 ± 0.045	20.789 ± 0.011
σ_{had} [nb]	41.541 ± 0.037	41.477 ± 0.009
R_b	0.21629 ± 0.00066	0.21576 ± 0.00004
R_c	0.1721 ± 0.0030	0.17227 ± 0.00004
A_{FB}^e	0.0145 ± 0.0025	0.01633 ± 0.00021
A_{FB}^μ	0.0169 ± 0.0013	0.01633 ± 0.00021
A_{FB}^τ	0.0188 ± 0.0017	0.01633 ± 0.00021
$A_\tau(P_\tau)$	0.1439 ± 0.0043	0.1475 ± 0.0010
$A_e(P_\tau)$	0.1498 ± 0.0049	0.1475 ± 0.0010
A_{FB}^b	0.0992 ± 0.0016	0.1034 ± 0.0007
A_{FB}^c	0.0707 ± 0.0035	0.0739 ± 0.0005
A_{LR}	0.15138 ± 0.00216	0.1475 ± 0.0010
m_W [GeV]	80.420 ± 0.031	80.381 ± 0.014
g_L^2	0.3009 ± 0.0028	0.3040 ± 0.0002
g_R^2	0.0328 ± 0.0030	0.03001 ± 0.00002
$g_V^{\nu e}$	-0.040 ± 0.015	-0.0398 ± 0.0003
$g_A^{\nu e}$	-0.507 ± 0.014	-0.5064 ± 0.0001
$Q_W(Cs)$	-73.20 ± 0.35	-73.23 ± 0.02

Table 4. Experimental values and SM predictions of the 21 different EWPO.

parameter, namely the dimensionless ratio v_{SM}/f with $v_{SM} = 246$ GeV, where f is the spontaneous symmetry breaking scale of the respective global symmetries. The ratio v_{SM}/f varies in the interval $[0,1]$: the lower bound is the *SM-* or *decoupling-limit* where all the modifications due to the Little Higgs structure are vanishing, recovering the SM results, while the upper limit is set in order to constrain the global symmetry breaking scale to be greater than the EWSB scale. The other free parameters are model-dependent, and under few assumptions we will consider only one extra free parameter for each of the three models.

In the L^2H model, the other free parameters are the mixing angles c, c' in the gauge sector (2.13), the ratio $R = \lambda_1/\lambda_2$ of the couplings in the top sector (2.43), and the parameter x proportional to the triplet *vev* (2.8). However, we will fix the value of R to a reference value of one, since our results will be with good approximation independent on the particular value of R, as a consequence of the collective symmetry breaking mechanism, as we will show later. For the remaining free parameters, we will let only the mixing angle c vary between $[0.1, 0.995]$, while presenting our results for few different choices of x and c'.

In the LHT model, besides the scale f, the only other free parameter for our study is again the ratio $R = \lambda_1/\lambda_2$ of the couplings in the T-even top sector (2.57). In Ref. [16]

the authors have performed a study to fix the allowed range for R in LHT: they obtained $R \lesssim 3.3$ by calculating the $J = 1$ partial-wave amplitudes in the coupled system of $(t\bar{t},$ $T\bar{T}_+, b\bar{b}, WW, Zh)$ states to estimate the tree-level unitarity limit of the corresponding scattering amplitudes. Therefore we will vary R between $[0.1, 3.3]$, where the lower limit is naïvely chosen by naturalness arguments.

In the SLH model the free parameters $R = \lambda_1^t/\lambda_2^t$ and μ_ϕ^2 are fixed by the requirement of EWSB, as described in Sec. 2 , leaving f and t_β as free parameters of our study, where t_β is the ratio of the vevs of the two scalar fields $\phi_{1,2}$ (2.24). We will let t_β vary between $[1.0, 15]$, where both limits are again naïvely chosen by naturalness. We will also require the perturbative constraint (2.35) to be satisfied: this will restrict the allowed values of t_β for a given value of f.

Following the procedure of Sec. 3, we determine the χ^2 measure (3.2) including the contributions from the deviations among the predicted and reported best-fit values of the different signal strength modifiers, and from the electroweak observables. The 95% and 99% CL allowed regions are then defined by the cumulative distribution function for an appropriate number of d.o.f.: having a total of 49 different best-fit channels, and 21 EWPO, the total number of d.o.f. is 70, since no free parameters have been fitted to the data.

We present in Fig. 1 the updated exclusion contours for the different models considered, distinguishing in particular Case A and Case B of the LHT. As the L^2H model is concerned, for the relative plot we have fixed $x = 0$ and $c' = 1/\sqrt{2}$, and restrict ourself to the region $v/f \in [0, 0.1]$ where the EWPO are satisfied and no tree-level decay of the Higgs involving new heavy partners is kinematically allowed, as already mentioned in Sec. 2. In the red region of the SLH plot the EWSB cannot be realized, and this region is therefore excluded. In the blue region of each plot we registered a total χ^2 which is lower than the χ^2 of the SM, with all signal strength modifiers set to 1 and all EWPO to their SM predictions. The minimum of the χ^2-value is denoted by the white points in the plots. The black lines represent contours of required fine-tuning inside the model setup, as we will explain later.

The new lower bounds of the symmetry breaking scale f at 99% and 95% CL within each model, as well as the value of f at which the minimum χ^2 has been determined, are summarized in Table 5.

	LHT **Case** A	LHT **Case** B	L^2H	SLH
$f_{\min}^{99\%}$ [TeV]	0.41	0.39	3.20	2.88
$f_{\min}^{95\%}$ [TeV]	0.47	0.39	3.58	3.26
$f_{\chi_{\min}^2}$ [TeV]	1.43	0.89	13.5	8.13
$\chi_{\min}^2$/d.o.f.	1.048	1.041	1.049	1.043
χ_{SM}^2/d.o.f.	1.054			

Table 5. 99% and 95% CL lower bounds on the symmetry breaking scale f and χ^2 comparison.

One should notice that there always exists a region in the parameter space where the measured χ^2 is equal or lower than the SM χ^2: however, the minimum χ^2 differs only at the 1% level w.r.t. the SM χ^2, so we can conclude that the agreement of these different
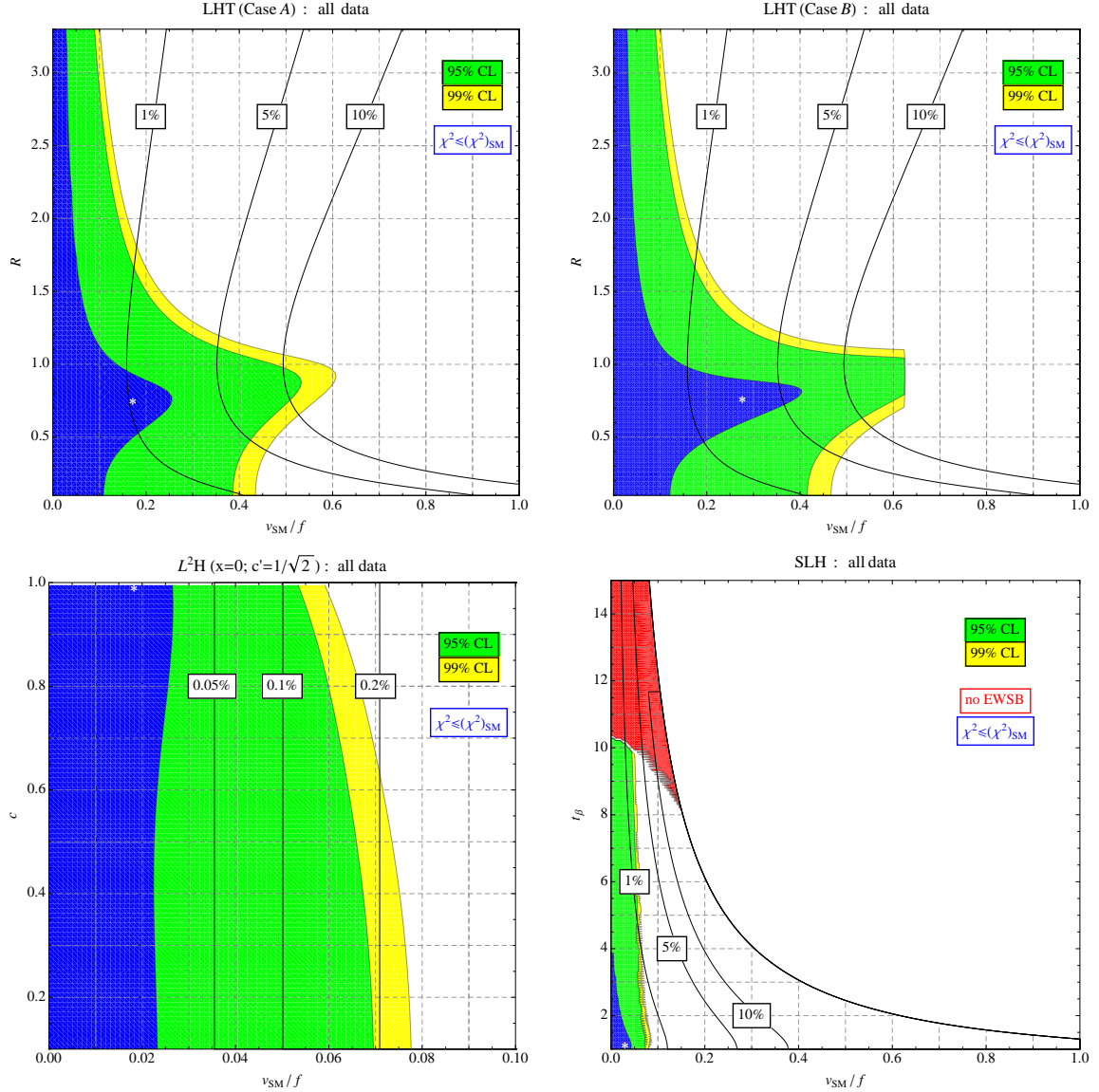
Figure 1. Allowed contours at 95% and 99% *CL* considering the whole available dataset for *LHT Case A* (up left), *LHT Case B* (up right), L^2H (down left) and *SLH* (down right). In the red region, no EWSB is possible. In the blue regions we found a lower χ^2 than the SM χ^2: the white points have the minimum χ^2. The thick black lines represent contours of required fine-tuning.

LH models with the collected data can be as good as within the SM, but not significantly better. In particular, for the L^2H and the *SLH* models the regions of equal or lower χ^2 than the SM χ^2 shrink to the SM-like decoupling limit.

The 99% *CL* lower limits of f can be translated into 99% *CL* lower limits on the spectrum of the new heavy particles of the different models, as summarized in Table 6.

The collective symmetry breaking mechanism implemented in each *LH* model eliminates all 1-loop quadratic divergences in the Higgs mass squared parameter, where the divergences from the SM particles are cancelled by quadratically divergent contributions

LHT	$m_{\min}$ [GeV], **Case A**	$m_{\min}$ [GeV], **Case B**
$m_{W_H} = m_{Z_H}$	269.6	262.2
m_{A_H}	64.5	62.8
m_Φ	291.7	283.7
m_{T_+}	553.6	537.5

L^2H	$m_{\min}$ [TeV]
$m_{W_H} = m_{Z_H}$	2.13
m_Φ	2.30
m_T	4.50

SLH	$m_{\min}$ [TeV]
m_{W_H}	1.35
m_{Z_H}	1.64
m_T	2.81

Table 6. 99% *CL* lower limits on the spectrum of the new heavy particles.

from new particles with same spin as the respective SM partners. The Higgs mass squared parameter is thus only logarithmic-divergent at 1-loop. As the masses of the new particles increase, the difference between the remaining SM 1-loop contributions and that of the new particles grows, requiring larger fine tuning of the Higgs mass squared parameter. The naturalness of the model could therefore be quantified observing by how much the contributions from the heavy states ($\delta\mu^2$) exceed the *observed* value of the Higgs mass squared parameter, as originally proposed in [5]:

$$\Delta = \frac{|\delta\mu^2|}{\mu^2_{\text{obs}}}, \qquad \mu^2_{\text{obs}} = \frac{m_h^2}{2}. \tag{4.2}$$

For example, if the new contributions to the Higgs mass squared parameter exceed μ^2_{obs} by a factor of 5, i.e. $\Delta = 5$, one says that the model requires 20% of fine tuning. Clearly, the lower the value of fine tuning, the worse is the naturalness of the model.

The dominant log-divergent contribution to the Higgs mass squared parameter comes from the top and its heavy partner loops, and is given for all the three *LH* models we are considering by [5]

$$\delta\mu^2 = -\frac{3\lambda_t^2 m_T^2}{8\pi^2} \log \frac{\Lambda^2}{m_T^2} \tag{4.3}$$

where $\Lambda = 4\pi f$ is the cut-off of the non-linear sigma model, λ_t is the SM top Yukawa coupling and m_T is the mass of the heavy top partner as defined in the different models. The thick black lines on the plots of Fig. 1 enclose these regions of required fine tuning (on the right hand side of the line), and the level of fine-tuning is also denoted on the plots.

We can see that the lowest level of fine-tuning is $\sim 10\%$ for both Case *A* and *B* of *LHT*, while significantly worse for both *SLH* ($\sim 1\%$) and L^2H ($\sim 0.1\%$). Comparing the naturalness of the model, accommodating the 7 and 8 TeV LHC results and the EWPO, to the MSSM [28] shows that only the model with *T*-parity, *LHT*, has less fine-tuning than the $\sim 1\%$ of the MSSM (in certain regions of parameter space). This is because the implemented *T*-parity relieves the constraints from EWPO, allowing a smaller value of the symmetry breaking scale f and therefore a smaller mass for the *T*-even top partner.

It is interesting to consider separately the χ^2 contributions from the best-fit values and from the EWPO. Considering e.g. the *LHT* case, the resulting plots are given in Fig. 2 [2].
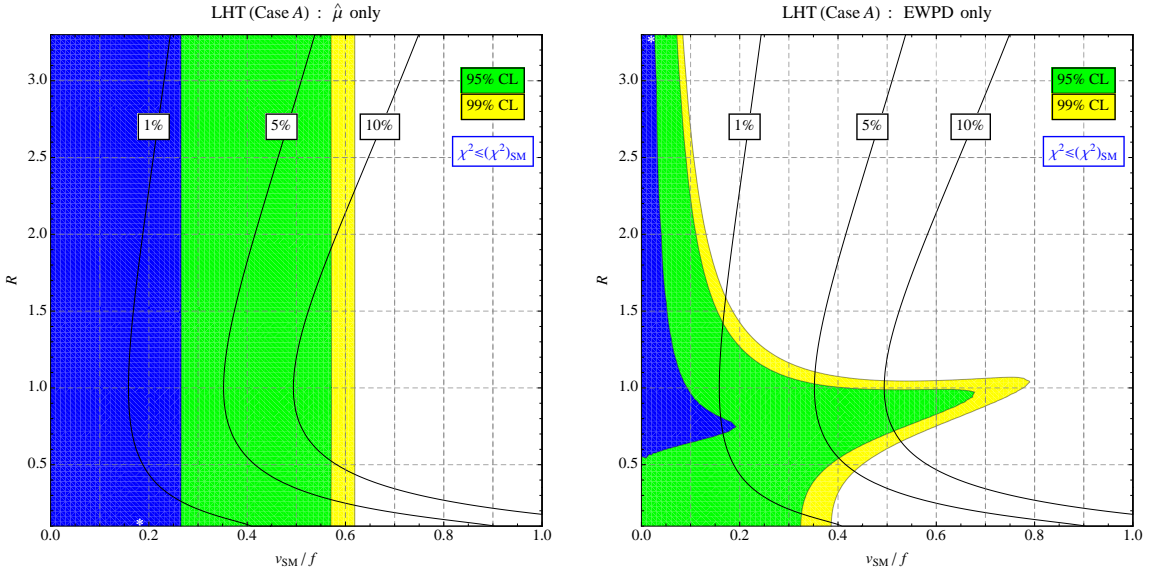
Figure 2. Allowed contours at 95% and 99% *CL* for *LHT Case A* considering separately the contributions from the Higgs sector only (left) and from EWPO only (right).

Clearly, the combined results Fig. 1 are mainly driven by the electroweak data: this should not be surprising, since the uncertainties on the electroweak observables are much smaller compared to the uncertainties on the best-fit values of the signal strength modifiers. This latter statement is therefore true also for the other *LH* models considered.

If only the Higgs data are considered, one can notice that there is a subdominant dependence on the parameter R compared to the ratio v_{SM}/f: this recovers a result already pointed out in the context of the Higgs Low-Energy Theorem (*LET*) in Composite Higgs models [51–55], namely that the effective *hgg* and *h$\gamma\gamma$* vertices do not depend on the details of the heavy fermion sector (in our case on R). We will give the argument in the sequel.

Focusing only on the partial decay widths into two gluons, in the *LET* approximation the interaction of the Higgs boson with gluons mediated by loops of colored particles, can be expressed (at leading order in the expansion of the Higgs field h around its vacuum expectation value v and considering only the contributions from fermions) by the following effective Lagrangian [55]

$$\mathcal{L}_{\text{hgg}} = \frac{g_s^2}{48\pi^2} G_{\mu\nu}^a G^{a\,\mu\nu} \frac{h}{v} \left[\frac{1}{2} v \frac{\partial}{\partial v} \log \det \mathcal{M}^\dagger \mathcal{M}_{|h=v} \right] , \qquad (4.4)$$

where $\mathcal{M}$ is the fermion mass matrix, including both the SM-like top and its heavy partner(s). In the narrow width approximation, the partial width into two gluons normalized to its SM value is given by the square of the expression in square brackets in Eq. (4.4), and agrees with our exact result (*A.6*) in the limit of heavy masses running in the loop.

---

[2]We consider only the results for Case *A*, as they are compatible with those of Case *B*

It is a general result [54] of Composite Higgs models, as well as Little Higgs models, that the determinant of the fermion mass matrix $\mathcal{M}^\dagger \mathcal{M}$ is only a function of the non-linear sigma model expansion parameter v/f and the details of the heavy fermion sector (i.e. on the masses and couplings of the fermions), but not separately on v:

$$\det \mathcal{M}^\dagger \mathcal{M}_{|h=v} = F\left(\frac{v}{f}\right) \times P\left(\lambda_i, m_i, f\right). \tag{4.5}$$

This factorization clearly makes both Eq. (4.4), and thus also the partial width into two gluons, independent of the couplings and masses of the fermions. It is only a function of the non-linear sigma model expansion parameter v/f.

Specializing to the LHT model case, one can easily see that this factorization indeed happens by considering the fermion mass matrix of Eq. (2.59):

$$\det \mathcal{M}^\dagger \mathcal{M}_{|h=v} = \frac{1}{2} \lambda_1^2 \lambda_2^2 f^4 \sin \frac{\sqrt{2}v}{f}$$
$$\frac{1}{2} v \frac{\partial}{\partial v} \log \det \mathcal{M}^\dagger \mathcal{M}_{|h=v} = 1 - \frac{2}{3}\frac{v}{f} \tag{4.6}$$

making the partial width into two gluons independent of the fermion couplings (i.e. of R), exactly in the LET limit and in good approximation with the exact expression, Eq. (A.6).

Analogous statements hold also for the effective coupling of the Higgs to two photons [55], so that we conclude that the partial width into two photons is exactly independent of R in the LET limit, and in good approximation with the exact expression, Eq. (A.7).

Note that the L^2H model shares with the LHT model the same form of the top-Yukawa Lagrangian, and therefore also the fermion mass matrix, cf. Eqs. (2.41) and (2.57). This allows us to take over the considerations about the dependence on model parameters for the L^2H model: in particular, the results of the Higgs sector will again not depend on the ratio R. Moreover, since the contribution from the heavy quark loop to the EWPO has been neglected (justified as in Ref. [15]), we were thus allowed to fix the value of R to a reference value (R=1) without loss of generality.

Another observation from the Higgs-only plot on the left of Fig. 2 is that the region with $v/f \gtrsim 0.62$ is highly disfavored by the collected data. This is due to the fact that for $f \lesssim 396.2$ GeV, the decay of the $m_h = 126$ GeV Higgs boson into two heavy photons A_H becomes open and dominant, highly reducing all other branching ratios and therefore the respective predicted signal strength modifiers μ^i (4.1), clearly in tension with the observed data, cf. Table 3.

An enhancement in the production cross sections could compensate this reduction in the branching ratios, but this is not the case for the LHT model, since all production modes are reduced w.r.t. their SM value. The VBF and HS production cross sections are slightly suppressed because of the suppressed coupling of the Higgs boson with SM gauge bosons, cf. Eqs. (A.11) and (A.12). The GF production cross section is also suppressed in the whole parameter space compared to the SM prediction. This suppression of the GF production in LHT was already pointed out in [14] in terms of the effective hgg coupling, Eq. (A.6). Indeed, the top Yukawa coupling is suppressed by means of the expansion of the non-linear

sigma model, Eq. (2.64), and the contribution from the T_+ partner (of opposite sign w.r.t. the top coupling) further suppresses the hgg coupling, as an effect of the collective symmetry breaking mechanism, as explained before in the context of the LET theorem, cf. Eq. (4.6). The contribution from the three T-odd partners u_1 is also negative, as we can see from Eq. (2.53). All these contributions cause thus a suppression in the effective hgg coupling (A.6) compared to its SM value.

Considering now the L^2H model, as for the LHT case the combined result on the lower left of Fig. 1 is mainly driven by the EWPO contribution, and even more dramatically since T-parity eliminates all tree-level contributions to the oblique parameters from the heavy states and from the triplet vev v', which are on the contrary present in the L^2H case.
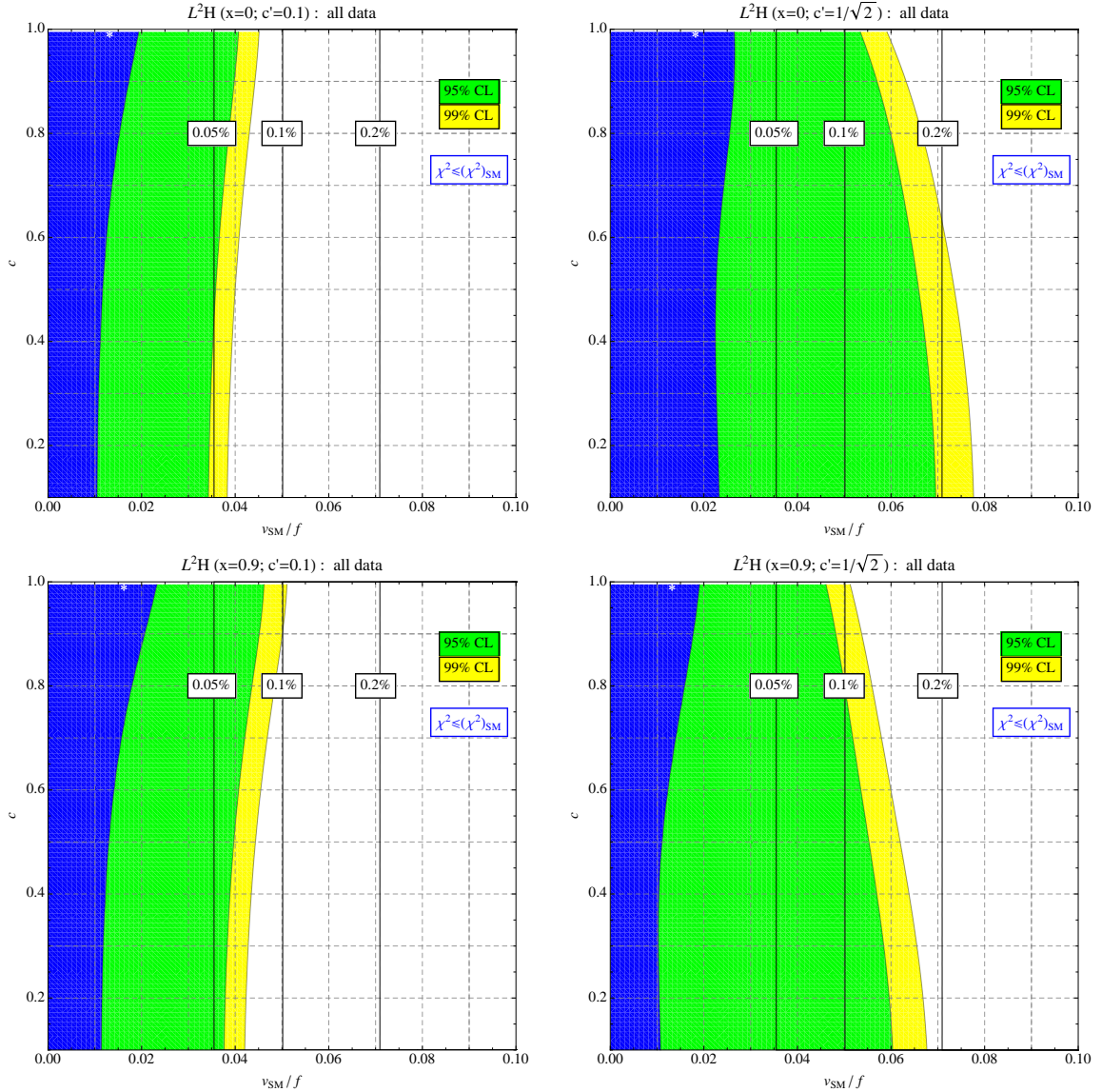
Figure 3. Allowed contours at 95% and 99% CL for L^2H with different choices of c' and x.

In Fig. 3, we show the dependence of the results on different choices of the parameters c' and x, in particular with $c' = \{0.1, 1/\sqrt{2}\}$, i.e. with minimal and maximal mixing, respectively, and with $x = \{0.0, 0.9\}$, i.e. with vanishing or nearly maximal triplet *vev*.

We can see that there is only a smooth dependence on c' and x: the result is driven mainly by the value of the symmetry breaking scale f. Compared to the *LHT* case, the parameter space of the L^2H model is indeed highly constrained for lower values of v/f, in particular $v/f \lesssim 0.1$, and this translates into a higher amount of required fine tuning, of the order $\sim 0.1\%$.
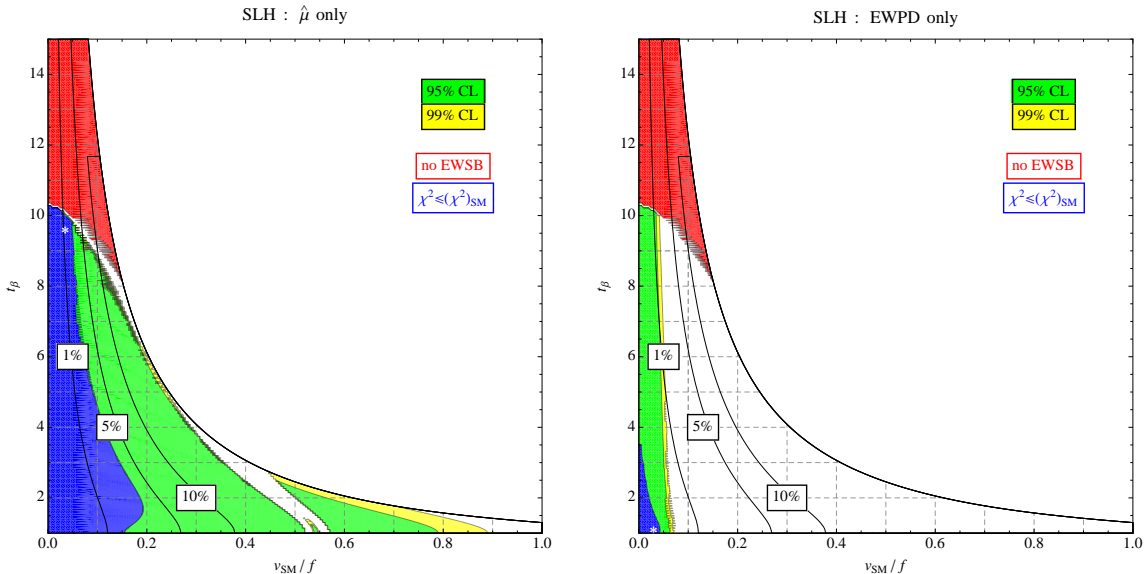
Figure 4. Allowed contours at 95% and 99% *CL* for *SLH* considering separately the contributions from the Higgs sector only (left) and from EWPO only (right).

Not surprisingly, also the results for the *SLH* are driven by the EWPO constraints, as can be seen from Fig. 4.

Nearly the whole parameter space is indeed still compatible at 99% *CL* with the Higgs-sector data: only new data with increasing luminosity and reduced uncertainties on the best-fit values of the signal-strength modifiers could give us more stringent information.

It is to be noted that in the Higgs-data plot there is however an excluded central band up to $v_{SM}/f \sim 0.6$: in this region, the decays of the Higgs involving the pseudo-scalar η are indeed open and dominant, cf. Eq. ($A.4$) and ($A.5$), highly reducing all other SM-like branching ratios, in the same way as for the decay of the Higgs into a pair of heavy photons A_H discussed in the *LHT* case. Indeed if we plot the ratio of the Higgs mass w.r.t. the mass of the pseudo-scalar η, we can identify the excluded regions in the left plot of Fig. 4 with the regions where the tree level decays involving the pseudo-scalar are kinematically accessible, cf. Fig. 5.
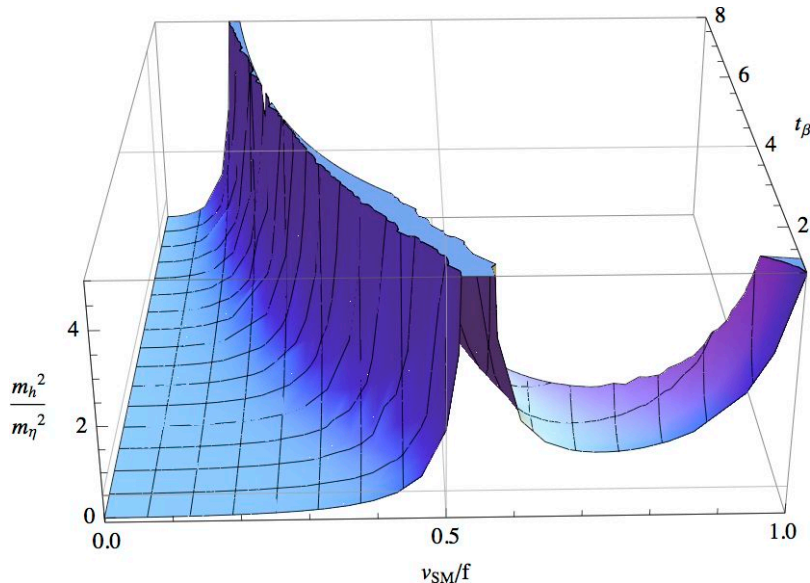
Figure 5. Higgs mass versus pseudo-scalar η mass in *SLH*.

5   Conclusions

In this paper, we have investigated the parameter space of the most prominent member of Little Higgs models, the Littlest Higgs model with and without T-parity from the class of Product Group Models, as well as the Simplest Little Higgs from the class of Simple Group Models, in the light of all present collider data as of the end of the year 2012. We included all published discovery and search channels for the Higgs boson from both LHC collaborations, ATLAS and CMS, together with the electroweak precision observables. The latter have been mostly updated by the results on the W mass from the Tevatron experiments recently.

Our results show that the experimental data from LHC are not yet precise enough to compete with the electroweak precision data, and we have to wait for un update from the collaborations with higher luminosity. For both the Littlest Higgs and the Simplest Little Higgs model without T-parity, EWPO force the Little Higgs scale f to be of the order of 2-4 TeV in order to be compatible with the precision electroweak data. Both models hence show a bit worse fine-tuning than the so-called natural pMSSM. On the other hand, T-parity does the job for which it has been invented, namely to reduce this amount of fine-tuning. The scale in the case of the Littlest Higgs model with T-parity is only constrained to be above 700-1200 GeV, and in most cases, the new T-odd particles can still be well below the TeV scale. The fine-tuning in the Higgs sector is less by a factor of two to five compared to the natural pMSSM.

We note further, that we did not include searches for exotic particles like additional gauge bosons or heavy vector-like quarks in our fits. This has been partially done elsewhere [28, 56, 57], and on the other hand, these searches will not be reaching enough sensitivity before the start of the 14 TeV run to become truly compatible.

Acknowledgments

A Higgs Boson Partial Widths and Production Cross Sections

For tree-level decays of the Higgs, the partial widths get a correction at lowest order via the corresponding modified couplings [14]:

$$\Gamma(h \to VV) = \Gamma(h \to VV)_{SM} \left(\frac{g_{hVV}}{g_{hVV}^{SM}} \right)^2 \qquad V \equiv W, Z$$

$$\Gamma(h \to \bar{f}f) = \Gamma(h \to \bar{f}f)_{SM} \left(\frac{g_{hff}}{g_{hff}^{SM}} \right)^2 \qquad f \equiv c, b, \mu, \tau \qquad \text{(A.1)}$$

with the different couplings and masses defined as in the previous sections.

There are also some new tree-level decay channels which are special to the different Little Higgs models, and which have to be taken into account if kinematically accessible. In particular, defining

$$x_i = \frac{4 \, m_i^2}{m_h^2} \qquad \text{(A.2)}$$

in LHT the Higgs field could decay into two heavy photons A_H, basically an invisible decay, with partial width [14]

$$\Gamma(h \to A_H A_H) = \frac{g_{hA_H A_H}^2 \, m_h^3}{128 \, \pi \, m_{A_H}^4} \sqrt{1 - x_{A_H}} \left(1 - x_{A_H} + \frac{3}{4} x_{A_H}^2 \right) \qquad \text{if } x_{A_H} < 1 \qquad \text{(A.3)}$$

while in SLH two new decay channels involving the pseudo-scalar η are possibly open [22]

$$\Gamma(h \to \eta\eta) = \frac{m_\eta^4}{8 \, \pi \, v^2 \, m_h} \sqrt{1 - x_\eta} \qquad \text{if } x_\eta < 1 \qquad \text{(A.4)}$$

$$\Gamma(h \to Z\eta) = \frac{m_h^3}{32 \, \pi \, f^2} \left(\frac{t_\beta^2 - 1}{t_\beta} \right)^2 \lambda^{3/2} \left(1, \frac{m_Z^2}{m_h^2}, \frac{m_\eta^2}{m_h^2} \right) \qquad \text{(A.5)}$$

with $\lambda(1, x, y) = (1 - x - y)^2 - 4xy$.

Defining the functions [25]

$$F_0(x) = x \left[1 - x f(x) \right]$$
$$F_{1/2}(x) = -2x \left[1 + (1 - x) f(x) \right]$$
$$F_1(x) = 2 + 3x + 3x(2 - x) f(x)$$
$$f(x) = \begin{cases} \left[\sin^{-1}\left(\frac{1}{\sqrt{x}} \right) \right]^2 & \text{for } x \geq 1 \\ -\frac{1}{4} \left[\log\left(\frac{1 + \sqrt{1-x}}{1 - \sqrt{1-x}} \right) - i\pi \right]^2 & \text{for } x < 1 \end{cases}$$

the general expression of the partial widths for the one-loop decays of the Higgs boson into two gluons or two photons are given by

$$\Gamma(h \to gg) = \frac{\alpha_s^2 \, m_h^3}{32 \, \pi^3 \, v^2} \left| \sum_{f, \, \text{col}} -\frac{1}{2} F_{1/2}(x_f) \, y_f \right|^2 \qquad , \qquad \text{(A.6)}$$

where the sum is extended over all colored fermionic particles of the spectrum which have a non-negligible coupling y_f to the Higgs boson (in contrast to SUSY there are no colored scalars in Little Higgs models), and by

$$\Gamma(h \to \gamma\gamma) = \frac{\alpha^2 \, m_h^3}{256 \, \pi^3 \, v^2} \left| \sum_{f, \text{ch}} \frac{4}{3} \, F_{1/2}(x_f) \, y_f + \sum_{v, \text{ch}} F_1(x_v) \, y_v + \sum_{s, \text{ch}} F_0(x_s) \, y_s \right|^2 \tag{A.7}$$

respectively. Here, the different sums run over all electrically charged fermionic (f), vector (v) and scalar (s) particles of the spectrum which have a non-negligible coupling $y_{f,v,s}$ to the Higgs boson.

At the LHC, the main production channel for the Higgs is the Gluon Fusion (GF):
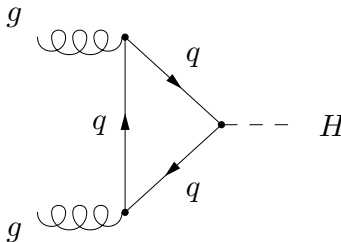
The hadronic GF cross section is given by the usual convolution [25]

$$\sigma \, (pp \to h) = \int_\tau^1 \frac{dx}{x} g \left(x, \mu_F^2 \right) g \left(\frac{\tau}{x}, \mu_F^2 \right) \hat{\sigma} \, (gg \to h) \, \tau \tag{A.8}$$

where $\tau = m_h^2/s$, with s the total hadronic c.m. energy squared, $g \left(x, \mu_F^2 \right)$ is the parton distribution function of the gluon at the factorization scale μ_F^2, and

$$\hat{\sigma} \, (gg \to h) = \frac{\pi^2}{8 \, m_h^3} \Gamma \, (h \to gg) \tag{A.9}$$

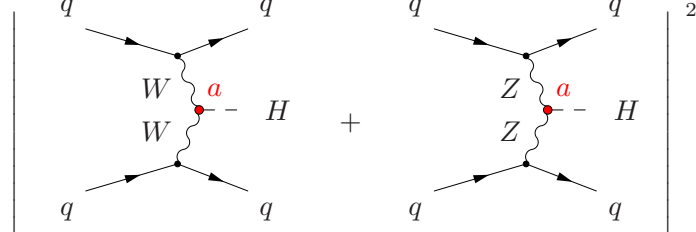is the partonic cross section in the narrow-width approximation.

Using Eq. ($A.9$) we can thus approximate the Little Higgs prediction for the GF cross section as a rescaling of the SM prediction:

$$\sigma(pp \to h)_{\text{LH}} \sim \frac{\Gamma \, (h \to gg)_{LH}}{\Gamma \, (h \to gg)_{SM}} \cdot \sigma(pp \to h)_{\text{SM}}. \tag{A.10}$$

The second most important channel for Higgs production at the LHC is the vector-boson Fusion (VBF). For the mass of the Higgs we are considering, the SM VBF cross section is smaller than the GF cross section by about an order of magnitude, but it could be important for some Higgs decay channels because of its distinctive kinematic signature of forward jets with high transverse momentum.

To calculate the Little Higgs prediction of the VBF cross section we have not included the contributions from the heavy gauge bosons (or any other heavy particles) as they are more difficult to be produced, and therefore we can safely neglect their contribution to the VBF cross section. We are left therefore with only the contributions from the light quarks $(u, \dots, b)$ and from the Z and $W^\pm$ gauge bosons.

Neglecting for simplicity possible $\mathcal{O}\left(v/f\right)$ corrections in the charged- and neutral-current couplings of the light quarks with the SM gauge bosons, at tree-level the Little Higgs VBF cross section is then given by its SM value rescaled with the appropriate (and model dependent) Higgs-gauge bosons coupling squared $a^2 \equiv (g_{hVV}/g_{hVV}^{SM})^2$. Hence, we can factorize the rescaling factor out of the amplitude
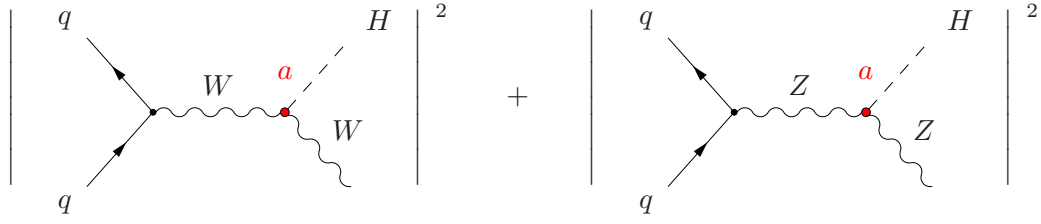

$$\Rightarrow \quad \sigma(qq \to qqh)_{\mathrm{LH}} = \left(\frac{g_{hVV}}{g_{hVV}^{SM}}\right)^2 \cdot \sigma(qq \to qqh)_{\mathrm{SM}}. \tag{A.11}$$

Even if the latter equation is a tree-level result, we have used the value of the SM VBF cross section recommended in [50], obtained at higher perturbative orders, to obtain the Little Higgs prediction through Eq. (A.11).

One should notice that in order to evaluate the VBF cross section in the SLH model, we neglected also the custodial-symmetry violating shift in the Z-mass (2.39), which is parametrically smaller than the other $\mathcal{O}\left(v^2/f^2\right)$ contribution, so that both hWW and hZZ vertices could share the same rescaling factor.

The Higgs-Strahlung production (HS) has a cross section which is about one to two orders of magnitude smaller than the GF cross section for Higgs masses $<$ 200 GeV, but it is important for certain Higgs decay channels because of the possibility of tagging the associated vector boson in leptonic decays. In particular one can distinguish between the production of the Higgs associated with the charged $W^{\pm}$ or with the neutral Z.

Using the same arguments as for the VBF case, we can conclude that at tree-level the correction to the HS cross section is again proportional to the square of the respective modified Higgs-gauge bosons coupling

$$\Rightarrow \quad \sigma(qq \to Vh)_{\mathrm{LH}} = \left(\frac{g_{hVV}}{g_{hVV}^{SM}}\right)^2 \cdot \sigma(qq \to Vh)_{\mathrm{SM}}. \tag{A.12}$$

As before, even if the latter equation is a tree-level result, we have used the value of the SM HS cross sections recommended in [50], obtained at higher perturbative orders, to obtain the Little Higgs prediction through Eq. ($A.12$).

B General structure of Electroweak Precision Observables

Using the notation of Ref. [32], one can parametrize the change of the charged- and neutral-current couplings of the SM-like gauge bosons due to the presence of new-physics as follows

$$\mathcal{L}_{cc} \supset -\frac{g}{\sqrt{2}} \sum_{i,j} \bar{f}_i \gamma^\mu \left((h_L + \delta\tilde{h}_L)P_L + (h_R + \delta\tilde{h}_R)P_R \right) f_j W_\mu$$

$$\mathcal{L}_{nc} \supset -\frac{g}{\tilde{c}_W} \sum_i \bar{f}_i \gamma^\mu \left((g_L + \delta\tilde{g}_L)P_L + (g_R + \delta\tilde{g}_R)P_R \right) f_i Z_\mu \quad , \tag{B.1}$$

where $g_{L,R}, h_{L,R}$ are normalized such that $g_L = I_3 - Q \cdot \tilde{s}_W^2$, $g_R = -Q \cdot \tilde{s}_W^2$, $h_L = V_{ij}$, $h_R = 0$. $P_{L,R}$ are the chiral projectors, and $\tilde{c}_W = g/\sqrt{g^2 + g'^2}$ (and similarly $\tilde{s}_W$) are the *bare* couplings as in the Standard Model[3].

In Ref. [32], the authors presented a parametrization of 21 different EWPO, both at low energies and at the Z pole, in terms of the oblique parameters S, T, U and of the coefficients $\delta\tilde{g}_{L,R}, \delta\tilde{h}_{L,R}$ defined as before. We refer to the original reference for the explicit expressions.

In the L^2H model, Ref. [15] already provides the explicit expressions of the 21 EWPO in terms of the free parameters of the model. For the LHT and SLH models, we calculated the different contributions of these models to the oblique parameters S, T, U and to the coefficients $\delta\tilde{g}_{L,R}, \delta\tilde{h}_{L,R}$, in order to reconstruct the different contributions to the EWPO, which have then been included in our χ^2 analysis.

---

[3]When evaluating electroweak observables, one should notice that the bare couplings $\tilde{c}_W, \tilde{s}_W, \tilde{e}$ can also get corrections from non-zero oblique parameters S,T,U.

References

[1] S. Chatrchyan *et al.* [CMS Collaboration], "Observation of a new boson at a mass of 125 GeV with the CMS experiment at the LHC," Phys. Lett. B **716**, 30 (2012) [arXiv:1207.7235 [hep-ex]].

[2] G. Aad *et al.* [ATLAS Collaboration], "Observation of a new particle in the search for the Standard Model Higgs boson with the ATLAS detector at the LHC," Phys. Lett. B **716**, 1 (2012) [arXiv:1207.7214 [hep-ex]].

[3] N. Arkani-Hamed, A. G. Cohen and H. Georgi, "Electroweak symmetry breaking from dimensional deconstruction," Phys. Lett. B **513**, 232 (2001) [hep-ph/0105239].

[4] N. Arkani-Hamed, A. G. Cohen, T. Gregoire and J. G. Wacker, "Phenomenology of electroweak symmetry breaking from theory space," JHEP **0208**, 020 (2002) [hep-ph/0202089].

[5] N. Arkani-Hamed, A. G. Cohen, E. Katz and A. E. Nelson, "The Littlest Higgs," JHEP **0207**, 034 (2002) [hep-ph/0206021].

[6] I. Low, W. Skiba and D. Tucker-Smith, "Little Higgses from an antisymmetric condensate," Phys. Rev. D **66**, 072001 (2002) [hep-ph/0207243].

[7] D. E. Kaplan and M. Schmaltz, "The Little Higgs from a simple group," JHEP **0310**, 039 (2003) [hep-ph/0302049].

[8] M. Schmaltz, "The Simplest little Higgs," JHEP **0408**, 056 (2004) [hep-ph/0407143].

[9] H. -C. Cheng and I. Low, "TeV symmetry and the little hierarchy problem," JHEP **0309**, 051 (2003) [hep-ph/0308199].

[10] M. Schmaltz and D. Tucker-Smith, "Little Higgs review," Ann. Rev. Nucl. Part. Sci. **55**, 229 (2005) [hep-ph/0502182].

[11] M. Perelstein, "Little Higgs models and their phenomenology," Prog. Part. Nucl. Phys. **58**, 247 (2007) [hep-ph/0512128].

[12] T. Han, H. E. Logan, B. McElrath and L. -T. Wang, "Phenomenology of the little Higgs model," Phys. Rev. D **67**, 095004 (2003) [hep-ph/0301040].

[13] S. R. Coleman and E. J. Weinberg, "Radiative Corrections as the Origin of Spontaneous Symmetry Breaking," Phys. Rev. D **7**, 1888 (1973).

[14] L. Wang and J. M. Yang, "The LHC di-photon Higgs signal predicted by little Higgs models," Phys. Rev. D **84**, 075024 (2011) [arXiv:1106.3916 [hep-ph]].

[15] C. Csaki, J. Hubisz, G. D. Kribs, P. Meade and J. Terning, "Big corrections from a little Higgs," Phys. Rev. D **67**, 115002 (2003) [hep-ph/0211124].

[16] A. Belyaev, C. -R. Chen, K. Tobe and C. -P. Yuan, "Phenomenology of littlest Higgs model with T^- parity: including effects of T^- odd fermions," Phys. Rev. D **74**, 115020 (2006) [hep-ph/0609179].

[17] J. Hubisz and P. Meade, "Phenomenology of the littlest Higgs with T-parity," Phys. Rev. D **71**, 035016 (2005) [hep-ph/0411264].

[18] J. Hubisz, P. Meade, A. Noble and M. Perelstein, "Electroweak precision constraints on the littlest Higgs model with T parity," JHEP **0601**, 135 (2006) [hep-ph/0506042].

[19] H. -C. Cheng and I. Low, "Little hierarchy, little Higgses, and a little symmetry," JHEP **0408**, 061 (2004) [hep-ph/0405243].

[20] C. -R. Chen, K. Tobe and C. -P. Yuan, "Higgs boson production and decay in little Higgs models with T-parity," Phys. Lett. B **640**, 263 (2006) [hep-ph/0602211].

[21] T. Han, H. E. Logan and L. -T. Wang, "Smoking-gun signatures of little Higgs models," JHEP **0601**, 099 (2006) [hep-ph/0506313].

[22] K. Cheung and J. Song, "Light pseudoscalar eta and H —> eta eta decay in the simplest little Higgs mode," Phys. Rev. D **76**, 035007 (2007) [hep-ph/0611294].

[23] F. del Aguila, J. I. Illana and M. D. Jenkins, "Lepton flavor violation in the Simplest Little Higgs model," JHEP **1103**, 080 (2011) [arXiv:1101.2936 [hep-ph]].

[24] W. Kilian, D. Rainwater and J. Reuter, "Distinguishing little-Higgs product and simple group models at the LHC and ILC," Phys. Rev. D **74**, 095003 (2006) [Erratum-ibid. D **74**, 099905 (2006)] [hep-ph/0609119].

[25] J. F. Gunion, H. E. Haber, G. L. Kane and S. Dawson, "The Higgs Hunter's Guide," Front. Phys. **80**, 1 (2000).

[26] W. Kilian and J. Reuter, "The Low-energy structure of little Higgs models," Phys. Rev. D **70**, 015004 (2004) [hep-ph/0311095].

[27] W. Kilian, D. Rainwater and J. Reuter, "Pseudo-axions in little Higgs models," Phys. Rev. D **71**, 015008 (2005) [hep-ph/0411213].

[28] J. Berger, J. Hubisz and M. Perelstein, "A Fermionic Top Partner: Naturalness and the LHC," JHEP **1207**, 016 (2012) [arXiv:1205.0013 [hep-ph]].

[29] M. Asano, S. Matsumoto, N. Okada and Y. Okada, "Cosmic positron signature from dark matter in the littlest Higgs model with T-parity," Phys. Rev. D **75**, 063506 (2007) [hep-ph/0602157].

[30] R. Contino, "The Higgs as a Composite Nambu-Goldstone Boson," arXiv:1005.4269 [hep-ph].

[31] R. Barbieri, B. Bellazzini, V. S. Rychkov and A. Varagnolo, "The Higgs boson from an extended symmetry," Phys. Rev. D **76**, 115008 (2007) [arXiv:0706.0432 [hep-ph]].

[32] C. P. Burgess, S. Godfrey, H. Konig, D. London and I. Maksymyk, "Model independent global constraints on new physics," Phys. Rev. D **49**, 6115 (1994) [hep-ph/9312291].

[33] G. Marandella, C. Schappacher and A. Strumia, "Little-Higgs corrections to precision data after LEP2," Phys. Rev. D **72**, 035014 (2005) [hep-ph/0502096].

[34] J. R. Espinosa, C. Grojean, M. Mühlleitner and M. Trott, "Fingerprinting Higgs Suspects at the LHC," JHEP **1205**, 097 (2012) [arXiv:1202.3697 [hep-ph]].

[35] J. R. Espinosa, C. Grojean, M. Mühlleitner and M. Trott, "First Glimpses at Higgs' face," JHEP **1212**, 045 (2012) [arXiv:1207.1717 [hep-ph]].

[36] A. Azatov, R. Contino and J. Galloway, "Model-Independent Bounds on a Light Higgs," JHEP **1204**, 127 (2012) [arXiv:1202.3415 [hep-ph]].

[37] A. Azatov and J. Galloway, "Electroweak Symmetry Breaking and the Higgs Boson: Confronting Theories at Colliders," arXiv:1212.1380 [hep-ph].

[38] G. Aad *et al.* [ATLAS Collaboration], "Observation of an excess of events in the search for the Standard Model Higgs boson in the γ-γ channel with the ATLAS detector", ATLAS-CONF-2012-091.

[39] G. Aad *et al.* [ATLAS Collaboration], "Observation and study of the Higgs boson candidate in the two photon decay channel with the ATLAS detector at the LHC", ATLAS-CONF-2012-168.

[40] G. Aad *et al.* [ATLAS Collaboration], "Search for the Standard Model Higgs boson in produced in association with a vector boson and decaying to bottom quarks with the ATLAS detector", ATLAS-CONF-2012-161.

[41] G. Aad *et al.* [ATLAS Collaboration], "Observation of an excess of events in the search for the Standard Model Higgs boson in the four lepton decay channel with the ATLAS detector", ATLAS-CONF-2012-169.

[42] G. Aad *et al.* [ATLAS Collaboration], "Updated ATLAS results on the signal strength of the Higgs-like boson for decays into WW and heavy fermion final states", ATLAS-CONF-2012-162.

[43] S. Chatrchyan *et al.* [CMS Collaboration], "Evidence for a new state decaying into two photons in the search for the standard model Higgs boson in pp collisions", CMS-PAS-HIG-12-015.

[44] S. Chatrchyan *et al.* [CMS Collaboration], "Observation of a new boson with a mass near 125 GeV", CMS-PAS-HIG-12-020.

[45] S. Chatrchyan *et al.* [CMS Collaboration], "Updated results on the new boson discovered in the search for the standard model Higgs boson in the $h \to ZZ \to 4l$ channel in pp collisions at $\sqrt{s} = 7$ and 8 TeV", CMS-PAS-HIG-12-041.

[46] S. Chatrchyan *et al.* [CMS Collaboration], "Evidence for a particle decaying to W^+W^- in the fully leptonic final state in a standard model Higgs boson search in pp collisions at the LHC", CMS-PAS-HIG-12-042.

[47] S. Chatrchyan *et al.* [CMS Collaboration], "Search for the standard model Higgs boson produced in association with W or Z bosons, and decaying to bottom quarks", CMS-PAS-HIG-12-044.

[48] S. Chatrchyan *et al.* [CMS Collaboration], "Serch for the standard model Higgs boson decaying to tau pairs", CMS-PAS-HIG-12-043.

[49] J. Beringer *et al.* [Particle Data Group Collaboration], "Review of Particle Physics (RPP)," Phys. Rev. D **86**, 010001 (2012).

[50] S. Dittmaier *et al.* [LHC Higgs Cross Section Working Group Collaboration], "Handbook of LHC Higgs Cross Sections: 1. Inclusive Observables," arXiv:1101.0593 [hep-ph]; https://twiki.cern.ch/twiki/bin/view/LHCPhysics/CrossSections

[51] A. Falkowski, "Pseudo-goldstone Higgs production via gluon fusion," Phys. Rev. D **77**, 055018 (2008) [arXiv:0711.0828 [hep-ph]].

[52] I. Low and A. Vichi, "On the production of a composite Higgs boson," Phys. Rev. D **84**, 045019 (2011) [arXiv:1010.2753 [hep-ph]].

[53] E. Furlan, "Gluon-fusion Higgs production at NNLO for a non-standard Higgs sector," JHEP **1110**, 115 (2011) [arXiv:1106.4024 [hep-ph]].

[54] A. Azatov and J. Galloway, "Light Custodians and Higgs Physics in Composite Models," Phys. Rev. D **85**, 055013 (2012) [arXiv:1110.5646 [hep-ph]].

[55] M. Gillioz, R. Grober, C. Grojean, M. Mühlleitner and E. Salvioni, "Higgs Low-Energy Theorem (and its corrections) in Composite Models," JHEP **1210**, 004 (2012) [arXiv:1206.7120 [hep-ph]].

[56] M. Perelstein and J. Shao, "T-Quarks at the Large Hadron Collider: 2010-12," Phys. Lett. B **704**, 510 (2011) [arXiv:1103.3014 [hep-ph]].

[57] S. Godfrey, T. Gregoire, P. Kalyniak, T. A. W. Martin and K. Moats, "Exploring the heavy quark sector of the Bestest Little Higgs model at the LHC," JHEP **1204**, 032 (2012) [arXiv:1201.1951 [hep-ph]].