Little Higgs models with a light T quark

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Abstract

We study Little Higgs models based on a $SU(3)_1 \times SU(3)_2$ global symmetry and with two scales (the two vacuum expectation values $f_{1,2}$) substantially different. We show that all the extra vector boson fields present in these models may be much heavier than the vectorlike $T$ quark necessary to cancel top-quark quadratic corrections. In this case the models become an extension of the standard model with a light ($\approx 500$ GeV) $T$ quark and a scalar Higgs field with a large singlet component. We obtain that the Yukawa and the gauge couplings of the Higgs are smaller than in the standard model, a fact that may reduce significantly the Higgs production rate through glu-glu and $WW$ fusion. The $T$-quark decay into Higgs boson becomes then a dominant Higgs production channel in hadron colliders.

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1 Introduction

The *naturalness* of the electroweak (EW) scale has been the main motivation to presume new physics below $\Lambda \approx 1 \text{ TeV}$. We have hints of a large grand unification scale, and we know that the Planck scale is there, so we *need* a mechanism that cancels the quantum corrections introduced by these large scales. The majority point of view has been that supersymmetry, technicolor or, more recently, extra dimensions could do the job and rise the cutoff of the standard model (SM) up to the fundamental scale.

This point of view, however, has become increasingly uneasy when facing the experimental evidence. Flavor physics, electric dipole moments and other precision electroweak observables suggest that, if present, the sfermion masses, the technicolor gauge bosons, or the Kaluza-Klein excitations of the standard gauge fields are above 5 TeV [1]. To be effective below the TeV and consistent with the data these models require a *per cent* fine tuning, whereas their presence at 5 TeV implies that nature deals with the Higgs mass parameter $m_h^2$ first using a mechanism to cancel 30 digits and then playing *hide and seek* with the last two digits. It may be more consistent either to presume that there should be another reason explaining this *little hierarchy* between the EW and the scale of new physics or that there is no dynamical mechanism that cancels *any* fine tuning in $m_h^2$ [2, 3]. This second possibility has been seriously considered after recent astrophysical and cosmological data suggested a non-zero vacuum energy density (the preferred value does not seem to be explained by any dynamics at that scale), and it will be clearly favored if no physics beyond the SM is observed at the LHC.

Little Higgs (LH) ideas [4, 5, 6] provide a very interesting framework with natural cancelations in the scalar sector. New symmetries protect the EW scale and define consistent models with a cutoff as high as $\Lambda \approx 10 \text{ TeV}$, scale where a more fundamental mechanism (SUSY [7] or extra dimensions [8]) would manifest. Therefore, these models could describe all the physics to be explored in the next generation of accelerators. More precisely, in LH models the scalar sector has a (tree-level) global symmetry that is broken spontaneously at a scale $f \approx 1 \text{ TeV}$. The SM Higgs doublet is then a Goldstone boson (GB) of the broken symmetry, and remains massless and with a flat potential at that scale. Yukawa and gauge interactions break explicitly the global symmetry. However, the models are built in such a way that the loop diagrams giving non-symmetric contributions must contain at least two different couplings. This *collective* breaking keeps the Higgs sector free of one-loop quadratic top-quark and gauge contributions (see [9, 10] for a recent review). Two types of models have been extensively considered in the literature: the ones based on a simple group (*littlest* $SU(5)$ model [5]) and the ones based on a product group (*simplest* $SU(3) \times SU(3)$ model...
of global symmetries.

Of course, an important point is then if the new physics that these models introduce is consistent with the data. In [12] it is shown that in general this is not the case, and the degree of fine tuning that they require is not below the one, for example, in the MSSM. LH models include an extra $T$ quark (that cancels quadratic top quark corrections) and massive gauge boson fields (that cancel quadratic gauge contributions and absorb extra GBs that otherwise would be massless). It is this later type of fields, the extra gauge bosons, the one giving large corrections to EW observables through mixing with the standard gauge bosons and through direct couplings with the light fermions. In models based on a global $SU(5)$ symmetry this problem can be solved imposing a $Z_2$ symmetry known as $T$ parity [13, 14]. This symmetry is analogous to the $R$ parity of SUSY models: under $T$ all the extra fields except for the $T$ quark are odd, whereas the standard fields are even. As a consequence, the mixing of standard and extra gauge bosons as well as the tree level exchange of extra bosons by standard fermions are forbidden. This keeps the corrections under control and allows $T$ quarks as low as 500 GeV, as required for an effective cancelation of quadratic corrections in the Higgs sector. Unfortunately, in the simplest $SU(3) \times SU(3)$ model it is not possible to implement a $T$ parity.

In this paper we explore the simplest LH model [11, 15, 16] and find that it can also accommodate a relatively light $T$ quark together with suppressed extra gauge boson contributions. This is achieved when the two VEVs $f_1$ and $f_2$ in these models are substantially different. We analyze this case and show that it has important phenomenological implications in Higgs physics. As we will argue, analogous results would be obtained in any LH model with a relatively low scale of global symmetry breaking.

2 The model

Let us start describing in some detail the model. The scalar sector contains two triplets, $\phi_1$ and $\phi_2$, of a global $SU(3)_1 \times SU(3)_2$ symmetry:

$$\phi_1 \rightarrow e^{i\theta_1 T^a} \phi_1, \quad \phi_2 \rightarrow e^{i\theta_2 T^a} \phi_2,$$

where $T^a$ are the generators of $SU(3)$. It is assumed that the scalar triplets get vacuum expectation values (VEVs) $f_1, f_2$ and break the global symmetry to $SU(2)_1 \times SU(2)_2$. The spectrum of scalar fields at this scale consists then of 10 massless modes (the GBs of the broken global symmetry) plus two massive fields (with masses of order $f_1$ and $f_2$). If one combination of the two global $SU(3)$ is made local, some of the GBs will be eaten by massive gauge bosons and the rest will define the SM Higgs sector.
In particular, if the two VEVs are

\[ \langle \phi_1 \rangle = \begin{pmatrix} 0 \\ 0 \\ f_1 \end{pmatrix}, \quad \langle \phi_2 \rangle = \begin{pmatrix} 0 \\ 0 \\ f_2 \end{pmatrix}, \quad (2) \]

and the diagonal combination of \( SU(3)_1 \times SU(3)_2 \) is made local,

\[ \phi_1 \to e^{i\theta_1 T^a} \phi_1, \quad \phi_2 \to e^{i\theta_2 T^a} \phi_2, \quad (3) \]

then the VEVs break \( SU(3) \times U(1)_\chi \) to the standard \( SU(2)_L \times U(1)_Y \), a process that takes 5 GBs. The other 5 GBs (the complex doublet \( (h^0, h^-) \) and a CP-odd singlet \( \eta \)) can be parametrized non-linearly \( [17] \):

\[ \phi_1 = e^{+i \frac{\tau_1}{f_1} \Theta} \begin{pmatrix} 0 \\ 0 \\ f_1 \end{pmatrix}, \quad \phi_2 = e^{-i \frac{\tau_2}{f_2} \Theta} \begin{pmatrix} 0 \\ 0 \\ f_2 \end{pmatrix}, \quad (4) \]

where

\[ \Theta = \frac{1}{f} \begin{pmatrix} \eta/\sqrt{2} & 0 & h^0 \\ 0 & \eta/\sqrt{2} & h^- \\ h^0 & h^- & \eta/\sqrt{2} \end{pmatrix}, \quad (5) \]

and \( f = \sqrt{f_1^2 + f_2^2} \).

If the global symmetry were exact, the Higgs boson would be massless. However, the symmetry is just \textit{approximate} (it is broken by top-quark and gauge-boson loops), so we expect that the non-symmetric operators will appear just suppressed by a loop factor. This may (should) give an acceptable VEV and a mass to the Higgs \( [18] \). Let us then assume that the real component of \( h^0 \) gets a VEV,

\[ \langle h^0 \rangle = u/\sqrt{2}, \quad (6) \]

and gives mass to the standard fermions and gauge bosons. This Higgs VEV implies the triplet VEVs

\[ \langle \phi_1 \rangle = \begin{pmatrix} if_1 s_1 \\ 0 \\ f_1 c_1 \end{pmatrix}, \quad \langle \phi_2 \rangle = \begin{pmatrix} -if_2 s_2 \\ 0 \\ f_2 c_2 \end{pmatrix}, \quad (7) \]

where

\[ s_1 \equiv \sin \frac{uf_2}{\sqrt{2}ff_1}, \quad s_2 \equiv \sin \frac{uf_1}{\sqrt{2}ff_2}. \quad (8) \]
To obtain the observed $W$ and $Z$ masses one needs

$$\sqrt{f_1^2 s_1^2 + f_2^2 s_2^2} = \frac{v}{\sqrt{2}} = 174 \text{ GeV}.$$  \hspace{1cm} (9)

Notice that using this non-linear realization of $\phi$, the Higgs VEV $u$ is not 246 GeV, although it goes to this value in the limit of large $f_1$ and $f_2$ (small $s_1$ and $s_2$).

In the unitary gauge all the GBs except for the singlet $\eta$ and the standard neutral Higgs $h$ are eaten by massive gauge bosons. In particular, it is easy to deduce the relation between $\phi_{1,2}$ and these fields:

$$\phi_1 = \exp \left( i \frac{f_2 \eta}{f_1 f \sqrt{2}} \right) \begin{pmatrix} i f_1 (s_1 \cos \frac{h_f}{\sqrt{2} f f_1} + c_1 \sin \frac{h_f}{\sqrt{2} f f_1}) & 0 \\ f_1 (c_1 \cos \frac{h_f}{\sqrt{2} f f_1} - s_1 \sin \frac{h_f}{\sqrt{2} f f_1}) & 0 \end{pmatrix},$$

$$\phi_2 = \exp \left( -i \frac{f_1 \eta}{f_2 f \sqrt{2}} \right) \begin{pmatrix} -i f_2 (s_2 \cos \frac{h_f}{\sqrt{2} f f_2} + c_2 \sin \frac{h_f}{\sqrt{2} f f_2}) & 0 \\ f_2 (c_2 \cos \frac{h_f}{\sqrt{2} f f_2} - s_2 \sin \frac{h_f}{\sqrt{2} f f_2}) & 0 \end{pmatrix}. \hspace{1cm} (10)$$

3 A light $T$ quark

As explained in the introduction, we are interested in models where the extra gauge bosons are heavy while the vectorlike $T$ quark that cancels quadratic corrections is lighter. The first requirement fixes the scale $f = \sqrt{f_1^2 + f_2^2}$, as the gauge boson masses are $\approx g f$. The top-quark Yukawa sector includes a triplet $\Psi_Q^T = (t \ b \ T)$ and two singlets $(t_1^c, t_2^c)$, and it is described by the Lagrangian

$$-\mathcal{L}_t = \lambda_1 \phi_1^\dagger \Psi_Q t_1^c + \lambda_2 \phi_2^\dagger \Psi_Q t_2^c + \text{h.c.},$$

where all the fermion fields are two-component spinors. In the limit of exact symmetry (i.e., $s_1 = 0 = s_2$) the top quark is massless and the extra $T$ quark has a mass $m_T$:

$$-\mathcal{L}_t \supset \lambda_1 f_1 T t_1^c + \lambda_2 f_2 T t_2^c + \text{h.c.} = m_T TT^c + \text{h.c.}, \hspace{1cm} (12)$$

where $m_T = \sqrt{\lambda_1^2 f_1^2 + \lambda_2^2 f_2^2}$ and $T^c = s_\alpha t_1^c + c_\alpha t_2^c$, with $s_\alpha = \lambda_1 f_1 / \sqrt{\lambda_1^2 f_1^2 + \lambda_2^2 f_2^2}$.

Once the Higgs VEV $u$ is included, it is easy to see that in order to have $m_T$ significantly smaller than $f$ we need $f_1 \ll f_2$ and $\lambda_2 \ll \lambda_1$. If we define $f_1 \equiv \epsilon f_2$, $\lambda_1 \equiv \lambda$ and $\lambda_2 \equiv \epsilon^2 \lambda$
this means that $\epsilon$ and $\epsilon'$ are small. At first order in these two parameters we have $f_2 \approx f$, $f_1 \approx \epsilon f$ and

$$
s_1 = \sin \frac{uf_2}{\sqrt{2}ff_1} \approx \sin \frac{u}{\sqrt{2}f},
$$
$$
s_2 = \sin \frac{uf_1}{\sqrt{2}ff_2} \approx \frac{\epsilon u}{\sqrt{2}f}.
$$

(13)

This implies (we redefine the top-quark field $-it \to t$)

$$
- L_t \supset \lambda \epsilon f \left(s_1 c_\alpha tt^c + s_1 s_\alpha tT^c + \frac{c_1}{s_\alpha} T T^c\right) + \text{h.c.},
$$

(14)

where $s_\alpha \equiv c_1 \epsilon / \sqrt{c_1^2 \epsilon^2 + \epsilon'^2}$ and

$$
t^c = c_\alpha t_1^c - s_\alpha t_2^c; \quad T^c = s_\alpha t_1^c + c_\alpha t_2^c.
$$

(15)

Taking $\epsilon f s_1 \approx v / \sqrt{2}$ we have

$$
- L_t \supset m_t tt^c + m_t t_\alpha tT^c + \frac{m_t}{c_\alpha s_\alpha t_1} T T^c + \text{h.c.},
$$

(16)

with $m_t \approx \lambda v c_\alpha / \sqrt{2}$, $t_\alpha = s_\alpha / c_\alpha$ and $t_1 = \tan(u / \sqrt{2} \epsilon f)$.

To obtain the mass eigenstates (we denote them through the paper with a prime) we still need to perform a rotation in the space of the left-handed fields $t$ and $T$:

$$
t' = c_\theta t - s_\theta T; \quad T' = s_\theta t + c_\theta T,
$$

(17)

which imply a heavy mass and a mixing

$$
m_T \approx \frac{m_t}{c_\alpha s_\alpha t_1}, \quad s_\theta \equiv V_{Tb} \approx \frac{m_t t_\alpha}{m_T}.
$$

(18)

### 4 Yukawa and gauge interactions

It is now easy to find the approximate Higgs couplings with the top and the $T$ quark. At the lowest order in $\epsilon$ and $\epsilon'$ we obtain

$$
- L_t \supset \frac{m_t}{v} \left(c_1 h t t^c + c_1 t_\alpha h T T^c - s_1 h t t^c - s_1 t_\alpha h T T^c\right) + \text{h.c.}
$$

(19)

for the Yukawas and

$$
- L_t \supset -\frac{1}{2m_T} \frac{m_t^2}{v^2} \left(s_1 c_1 h^2 t t^c + s_1 c_1 h^2 c_\alpha h T T^c + \frac{c_2}{s_\alpha c_\alpha} h^2 T T^c + \frac{c_2}{c_\alpha} h^2 T T^c\right) + \text{h.c.}
$$

(20)
This lagrangian exhibits two features. The first one is common to all LH models, namely, the quadratic corrections from the diagrams in Fig. 1 cancel (the correction that these diagrams introduce is logarithmic and proportional to $m_T^2$). The second one is that the top-quark Yukawa coupling $y_t$ is not $\sqrt{2}m_t/v$, like in the SM. Here the coupling appears suppressed by a factor of $c_1$. This is a generic feature in all LH models, and it can be understood as the contribution of higher dimensional operators in the non-linear expansion or as a mixing of order $v/\sqrt{2}f_1$ of the Higgs doublet with the $SU(2)_L$ singlets breaking the global symmetries. In the model under study the scale $f_1$ where $SU(3)_1$ is broken is relatively low, so the effect becomes important. Notice that only the doublet component of the Higgs couples to the fermions and gives them a mass $y_f^SMv/\sqrt{2}$. Then, if the doublet is just a component $c_1$ along the physical Higgs, the Yukawa couplings will be $y_f = y_f^SMc_1$. Actually, in the model under study we still have to perform the rotation in Eq. (17) to obtain the quark mass eigenstates. The flavor-diagonal Yukawa couplings are then

$$-\mathcal{L}_t \supset \frac{m_t}{v} \left( (c_1 c_\theta + s_1 s_\theta) h t' t'^c - (s_1 c_\alpha c_\theta - c_1 s_\alpha s_\theta) h T'^T t'^c \right) + \text{h.c.}, \quad (21)$$

which imply

$$\frac{y_t}{y_t^SM} \approx c_1 c_\theta + s_1 s_\theta \approx c_1 + s_1 V_{Tb}. \quad (22)$$

In the gauge sector we obtain the same type of suppression effect. The gauge couplings of the Higgs $h$ with both the $W$ and the $Z$ vector bosons are reduced with respect to the SM values by a factor of

$$\frac{g}{g^{SM}} = \frac{\sqrt{2}f_1 f_2 (s_1 c_1 + s_2 c_2)}{v f} \approx c_1, \quad (23)$$

where $g$ stands for the $SU(2)_L$ and $U(1)_Y$ couplings. In this model the singlet component $s_1$ along the Higgs, the mass $m_T$ of the vectorlike quark, and the mixing $V_{Tb}$ between $T$ and $t$ depend only on two parameters ($\epsilon$ and $\epsilon'$), which yields the approximate relation

$$\frac{s_1}{c_1} \approx V_{Tb} + \frac{m_T^2}{m_T^2 V_{Tb}}. \quad (24)$$
In Fig. 2 we plot the exact (numerical) correlation among these three quantities.

The approximate Yukawa couplings with the (pseudo) GB \( \eta \) can also be obtained expanding \( \phi_{1,2} \) in Eq. (10):

\[- \mathcal{L}_t \supset i \frac{m_t}{v} (s_1 \eta t t^c + s_1 t_a \eta t T^c + c_1 \eta T t^c + c_1 t_a \eta T T^c) + h.c. \] (25)

It is also easy to deduce the couplings of the top and the \( T' \) quarks with the gauge bosons. After the rotation in Eq. (17) and neglecting the mixings with the lighter quarks, the couplings with the \( W \) boson are

\[ \mathcal{L}_W = - \frac{g}{\sqrt{2}} \bar{t} \sigma_{\mu b} W_{\mu}^+ + h.c. \]

\[ = - \frac{g}{\sqrt{2}} \left( \sqrt{1 - V_{Tb}^2} \bar{T'} \sigma_{\mu b} + V_{Tb} T' \sigma_{\mu b} \right) W_{\mu}^+ + h.c. \] (26)

In the \( Z \)-boson sector we find additional flavour-changing interactions:

\[ \mathcal{L}_Z \supset - \frac{g}{2 c_W} \left( \bar{t} \begin{pmatrix} 1 & 0 \\ 0 & 0 \end{pmatrix} t \right) \sigma_{\mu} \left( \begin{pmatrix} t \\ T \end{pmatrix} \right) Z_{\mu} \]

\[ \approx - \frac{g}{2 c_W} \left( \bar{t}' \begin{pmatrix} 1 & 0 \\ V_{Tb}^2 & V_{Tb} \end{pmatrix} t' \right) \sigma_{\mu} \left( \begin{pmatrix} t' \\ T' \end{pmatrix} \right) Z_{\mu} \] (27)
coupling with the Z boson is suppressed by a factor of $c$.

At 500 GeV that cancels top quark corrections, a light quark plus a state of mass $m_T = 574$ GeV. The mass eigenstates are

$$t' = 0.98 t - 0.20 T ; \quad t'^* = 0.81 t^*_1 - 0.57 t^*_2 ,$$

$$T' = 0.20 t + 0.98 T ; \quad T'^* = 0.57 t^*_1 + 0.81 t^*_2 ,$$

which imply $V_{Tb} = 0.20$. The dimension four and five couplings of these fields with the Higgs read

$$- \mathcal{L}_t \supset \frac{m_t}{v} \left( 0.89 h t' t'^* - 0.62 h t' T'^* - 0.38 h T' t'^* - 0.27 h T' T'^* \right) + \text{h.c.}$$

and

$$- \mathcal{L}_t \supset - \frac{1}{2M_T v^2} \left( 0.73 h^2 t' t'^* + 0.51 h^2 t' T'^* + 1.69 h^2 T' t'^* + 1.19 h^2 T' T'^* \right) + \text{h.c.}$$

This means that $y_t/y_t^{SM} = 0.89$. The Yukawa couplings with $\eta$ are

$$- \mathcal{L}_t \supset i \frac{m_t}{v} \left( 0.38 \eta t' t'^* + 0.27 \eta t' T'^* + 0.89 \eta T' t'^* + 0.62 \eta T' T'^* \right) + \text{h.c.} ,$$

whereas the couplings of these quarks with the $W$ boson are

$$\mathcal{L}_W = - \frac{g}{\sqrt{2}} \left( 0.98 \bar{T} b + 0.20 \bar{T} b \right) \gamma^\mu W_\mu^+ + \text{h.c.}$$

The interactions with the $Z$ boson include the flavour-changing terms

$$\mathcal{L}_Z \supset - \frac{g}{2c_W} \left( \bar{T'} T' \right) \sigma^\mu \begin{pmatrix} 0.96 & 0.20 \\ 0.20 & 0.04 \end{pmatrix} \begin{pmatrix} t' \\ T' \end{pmatrix} Z_\mu ,$$

It is remarkable that in models with just a vectorlike $T$ quark the mass eigenstates couple to the Higgs only through the term that also introduces the mixing with the top quark, i.e., with a coupling $V_{Tb} m_T/v$ [19]. Here, however, the Higgs couples to $T$ and $T'^*$ even if the mixing $V_{Tb}$ is zero.

Finally, several comments about the stability of the scales are in order. First, notice that the natural cutoff of this model is at $\Lambda \approx 4 \pi f_1$. In the limit of $f_1 = v/\sqrt{2}$ (i.e., $s_1 = 1$) this is just the SM cutoff, whereas values of $f_1$ around 300 GeV rise the cutoff by a factor of 2 up to $\Lambda \approx 4$ TeV. Second, in order to decouple the extra gauge bosons we are taking a large $f_2$ scale, around the cutoff $\Lambda$. This defines a minimal LH model with only a $T$ quark at 500 GeV that cancels top quark corrections, a light singlet $\eta$, and a light Higgs $h$ whose coupling with the $Z$ boson is suppressed by a factor of $c_1$. This suppression should relax LEP bounds on its mass, which tend to be below 100 GeV.

\footnote{Being a gauge singlet, $\eta$ avoids LEP bounds.}
5 Electroweak precision observables

Let us start analyzing the implications on EW precision observables. The three basic ingredients of the LH models under study are the presence of heavy vector bosons, of a relatively light $T$ quark, and of a sizeable singlet component in the Higgs field.

(i) The massive gauge bosons would introduce mixing with the standard bosons and four fermion operators. This could manifest as a shift in the $Z$ mass and corrections in atomic parity violation experiments and LEP II data. However, none of these effects is observable if $f_2 \geq 3$ TeV [11].

(ii) The effects on EW precision observables due to the singlet component of the Higgs field are also negligible. Although the Yukawa coupling of the top with the neutral Higgs is here smaller than in the SM, it is the coupling with the would be GBs (the scalars eaten by the $W$ and $Z$ bosons) what determines the large top-quark radiative corrections, and these are not affected by the presence of singlets.

(iii) The bounds on a vectorlike $T$ quark from precision EW data have been extensively studied in the literature, we will comment here the results in [19] as they apply to LH models in a straightforward way.

The mixing of the top quark with the $T$ singlet reduces its coupling with the $Z$ boson. This, in turn, affects the top quark radiative corrections (triangle diagrams) to the $Zbb$ vertex, which is measured in the partial width $Z \to b\bar{b}$ \(R_b = \Gamma(Z \to b\bar{b})/\Gamma(Z \to \text{hadrons})\) and forward-backward asymmetries. The heavier $T$ quark also gives this type of corrections to the $Zbb$ vertex, and for low values of $m_T$ both effects tend to cancell (i.e., if $m_T = m_t$ the vertex $Zbb$ is the same as in the SM). For large values of $m_T$ (above 500 GeV) the upper bound on $V_{Tb}$ from precision $b$ physics is around 0.2 [19].

The $T$ quark would also appear in vacuum polarization diagrams, affecting the oblique parameters $S$, $T$, and $U$. For degenerate masses ($m_T = m_t$) the corrections to $T$ and $U$ vanish for any value of the mixing $V_{Tb}$ and the correction to $S$ is small ($\Delta S \approx -0.16V_{Tb}$). For large values of $m_T$ the only oblique parameter with a sizeable correction is $T$ ($\Delta T \approx 2.7V_{Tb}$ for $m_T = 500$ GeV), but the limits on $V_{Tb}$ are in this case smaller than the ones from $R_b$ [19].

6 Higgs physics

The phenomenological impact of these models on Higgs physics at hadron colliders may be important. The main effects can be summarized as follows.
Figure 3: Ratios $R_{gg} \equiv \sigma(gg \to h)/\sigma^{SM}(gg \to h)$ (solid) and $R_{WW} \equiv \sigma(WW \to h)/\sigma^{SM}(WW \to h)$ (dashes) for $V_{Tb} = 0.20, 0.15$ and different values of $m_T$. $R_{gg}$ and $R_{WW}$ coincide at the 1% level.

(i) Suppression of the $gg \to h$ cross section. This effect is due to the suppression of the top Yukawa coupling relative to the SM value (see Fig. 2) and also to the contribution of the extra $T$ quark. Although this second factor is numerically less important, it is remarkable that always interferes destructively in the amplitude: the relative minus sign versus the top-quark contribution follows from the cancelation of quadratic corrections to $m_T^2$. Notice that the two diagrams for $gg \to h$ are obtained from the diagrams in Fig. 1 just by adding two gluons to the fermion loop and changing a Higgs leg by its VEV.

It is easy to obtain approximate expressions for this suppression factor in the limit of $m_H \ll m_t, m_T$ \[20\]:

$$\frac{\sigma(gg \to h)}{\sigma^{SM}(gg \to h)} \approx \left(\frac{y_t}{v^{SM}} + \frac{y_{T\nu}}{m_T}\right)^2$$
$$\approx \left(c_1 c_\theta + s_1 s_\theta - t_1 s_\alpha (s_1 c_\theta - c_1 s_\theta)\right)^2$$
$$\approx c_i^2$$

In Fig. 3 we plot the ratio $R_{gg} \equiv \sigma(gg \to h)/\sigma^{SM}(gg \to h)$ for different values of $V_{Tb}$ and $m_T$. For $m_H = 150$ GeV the approximation above is good at the 1% level. This effect, which
could hide the Higgs at the LHC, has been recently discussed in general models with scalar singlets [21, 22] and also in the framework of LH models with T parity [23].

(ii) Suppression in the production cross sections that involve gauge interactions: $WW \to h$, $q\bar{q} \to Wh$, etc. [see Eq. (23)]. We plot the ratio $R_{WW} \equiv \sigma(WW \to h) / \sigma^{SM}(WW \to h) \approx c_1^2$ also in Fig 3. It is remarkable that for $m_H = 150$ GeV the suppression in these cross sections coincides with the one in $\sigma(gg \to h)$ at the 1% level.

(iii) New production channels through T-quark decay [24]. A T quark of mass below 600 GeV will be copiously produced at the LHC. In particular, the cross section to produce $TT$ pairs in $pp$ collisions goes from $10^4$ fb for $m_T = 400$ GeV to $10^3$ fb for $m_T = 600$ GeV [25]. Once produced, a T quark may decay into $Wb, Zt, ht,$ and $\eta t$ [26]. We find an approximate relation among the partial widths in the limit of $m_T$ much larger than the mass of the final particles:

$$\Gamma(T \to Wb) \approx \frac{\alpha}{16\alpha^2} \frac{V_{Tb}^2}{M_W} \frac{m_T^2}{M_W^2}$$

$$\Gamma(T \to Zt) \approx \frac{1}{2} \Gamma(T \to Wb)$$

$$\Gamma(T \to ht) \approx \frac{1}{2} \left( c_1^2 + \frac{s_1^2}{t_\alpha^2} \right) \Gamma(T \to Wb)$$

$$\Gamma(T \to \eta t) \approx \frac{1}{2} \left( s_1^2 + \frac{c_1^2}{t_\alpha^2} \right) \Gamma(T \to Wb)$$ (35)

Notice that the T quark will decay through the 4 channels with branching ratios that are independent of $V_{Tb}$. $T \to W^+b$ gives the best discovery potential for the T quark, whereas the Higgs $h$ will be produced with a branching ratio close to the 20%. The detailed signal and background study at the LHC in [25] shows that $TT \to W^+b\bar{b}h \to W^+bW^-\bar{b}h$ and $TT \to hth\bar{t} \to W^+bW^-\bar{b}h$ give a very high statistical significance for the Higgs (around 10σ for 30 fb$^{-1}$). We expect similar results in this model, although the presence of the scalar $\eta$ can open new decay channels for the Higgs. In particular, if $m_H > 2m_\eta$ the (global symmetry-breaking) coupling $h\eta\eta$ could loosen LEP bounds on the Higgs mass and open the interesting channel $h \to \eta\eta \to 4b$ [16]. In particular, if $m_H > 2m_\eta$ the coupling $h\eta\eta$ opens the interesting channel $h \to \eta\eta \to 4b$ [16] that, together with the suppression in the $hZZ$ coupling, could loosen considerably the 114 GeV LEP bound [27] on the Higgs mass.

7 Summary and discussion

LH models are minimal extensions of the SM that rise its natural cutoff. All these models contain a vectorlike T quark that cancels one-loop quadratic corrections to $m_h^2$. An effective
cancellation requires $m_T \approx 500 \text{ GeV}$, which implies a scale of global symmetry breaking of the same order. We have studied models based on a $SU(3)_1 \times SU(3)_2$ global symmetry and have shown that all the other ingredients of the models (namely, the extra vector and scalar fields) can be decoupled. In LH models based on a simple group this decoupling effect is achieved using a discrete symmetry known as $T$ parity, whereas here it is obtained fixing one of the VEVs ($f_2$) around 4 TeV and making the other one ($f_1$) up to a factor of $4\pi$ smaller.

The Higgs $h$ has then suppressed gauge and Yukawa couplings. This effect can be understood as a mixing of order $v/\sqrt{2}f_1$ of the doublet with the singlet ($\sigma$) that breaks the global symmetries and gives mass to the $T$ quark. This seems to be a generic feature in any LH models: the lighter is the extra $T$ quark that cancels top-quark corrections, the larger is the singlet component $s_1$ along the Higgs $h$. In our model the scalar $\sigma$ that mixes with the doublet gets a mass of order $f_1$. Notice that this scalar is necessary to unitarize the theory, since the gauge coupling of the light Higgs $h$ is here suppressed by a factor of $c_1$ (if its mass is around the cutoff $4\pi f_1$, in the limit $s_1 = 1$ the model becomes Higgsless).

The reduction in gauge and Yukawa couplings respect to the SM values have consequences at hadron colliders. In particular, the Higgs production rate through $gg$ and $WW$ fusion will be suppressed by a factor of $c_2^2$. These models are, actually, a realization of the ideas discussed in [28], where the Higgs field is spread into several weaker modes. It is obvious that this, together with the possible new decay mode $h \rightarrow \eta \eta$, could loosen substantially present bounds on the Higgs mass.

Although the standard channels to produce Higgs bosons are suppressed, the presence of a relatively light $T$ quark opens new possibilities. These fields will be copiously produced through tree-level interactions in hadron colliders, and they may decay into Higgs plus a top quark. General analysis that can be found in the literature do not consider the effect of a scalar singlet, as the mass of the quark is introduced ad hoc and not through scalar VEVs. We have taken the singlet into account and have shown that this decay mode has an approximate branching ratio of the 20% (versus the 25% in models with no singlets).

It is amusing that in these LH models the discovery of the Higgs at the LHC comes together with the discovery of a vectorlike quark and, thus, of a new scale in particle physics (notice that $T$ is not at the EW scale). If that were the case, the new natural cutoff of the model would be rised up to energies just above the reach of the LHC, which would certainly provide for good arguments to build a bigger collider.
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