

**This is an electronic reprint of the original article.
This reprint *may differ* from the original in pagination and typographic detail.**

Author(s): Antola, Matti; Di Chiara, Stefano; Tuominen, Kimmo

Title: Ultraviolet complete technicolor and Higgs physics at LHC

Year: 2015

Version:

Please cite the original version:

Antola, M., Di Chiara, S., & Tuominen, K. (2015). Ultraviolet complete technicolor and Higgs physics at LHC. *Nuclear Physics B*, 899, 55-77.
<https://doi.org/10.1016/j.nuclphysb.2015.07.012>

All material supplied via JYX is protected by copyright and other intellectual property rights, and duplication or sale of all or part of any of the repository collections is not permitted, except that material may be duplicated by you for your research use or educational purposes in electronic or print form. You must obtain permission for any other use. Electronic or print copies may not be offered, whether for sale or otherwise to anyone who is not an authorised user.



Ultraviolet complete technicolor and Higgs physics at LHC

Matti Antola ^{a,1}, Stefano Di Chiara ^{b,a,*}, Kimmo Tuominen ^{c,a}

^a *Helsinki Institute of Physics, P.O. Box 64, FI-000140, Univ. of Helsinki, Finland*

^b *Department of Physics, P.O. Box 35, FI-40014, Univ. of Jyväskylä, Finland*

^c *Department of Physics, P.O. Box 64, FI-000140, Univ. of Helsinki, Finland*

Received 23 February 2015; received in revised form 27 June 2015; accepted 11 July 2015

Available online 23 July 2015

Editor: Hong-Jian He

Abstract

We consider a supersymmetric model with a new strong interacting sector. The model is built on a strongly interacting $\mathcal{N} = 4$ Super Yang Mills sector, broken explicitly to $\mathcal{N} = 1$ supersymmetry by embedding within the Minimal Supersymmetric Standard Model (MSSM). Due to cancellation of global and gauge anomalies, the model additionally features a fourth lepton superfamily. We propose a scenario where all elementary scalars, gauging and higgsinos are decoupled at an energy scale substantially higher than the electroweak (EW) scale, thereby avoiding the little hierarchy problem of MSSM.

We construct a low energy effective model, where EW symmetry breaking and viable mass spectrum are produced dynamically. To test further the viability of the model, we work out the Higgs couplings as well as the EW precision parameters and then perform a goodness of fit analysis using LHC and EW precision data. The model fits the given experimental data at a level comparable to that of the Standard Model.

© 2015 The Authors. Published by Elsevier B.V. This is an open access article under the CC BY license (<http://creativecommons.org/licenses/by/4.0/>). Funded by SCOAP³.

* Corresponding author.

E-mail addresses: antola.matti@gmail.com (M. Antola), dichiara@helsinki.fi (S. Di Chiara), kimmo.i.tuominen@jyu.fi (K. Tuominen).

¹ Currently at Eniram, Helsinki, Finland.

1. Introduction

The data collected at the LHC experiments during the 7 and 8 TeV runs, with the epochal discovery of the Higgs boson [1,2] and the measurement of its couplings [3,4], seem to have provided the experimental verification of the Standard Model (SM) in its entirety below a TeV. Because of this, new physics coupled with the Electroweak (EW) SM currents must have a typical scale of the order of a TeV or higher. One possibility is that the new physics scale is much above the terascale, and then naturality as a model building paradigm should be reinterpreted [5,6]. If on the other hand naturality is realized in nature, then the new physics scale should be near the terascale and the present LHC data would provide hints of a new spectrum awaiting discovery in the future runs at the LHC. In this paper we investigate a model framework falling into the latter category.

In Technicolor (TC) theories [7,8] the new physics scale is naturally of the order of a TeV. The EW symmetry is broken by a new strong interaction which generates a fermion condensate, in a way analogous to QCD. The absence of light elementary scalars in TC automatically solves the SM fine-tuning problem. The mass of the lightest composite scalar is usually expected to be of $\mathcal{O}(\text{TeV})$, well above the measured 126 GeV value. A light scalar can arise as consequence of approximate global symmetries, chiral symmetry [9–11] or scale invariance [12–16]. However, only recently it has been realized that also with simple QCD-like TC dynamics the scalar particle can become light because of loop corrections originating from extended sectors, which are always required in TC models to account for the generation of fermion masses [17–19].

The observed mass pattern of matter fields is generated in TC by coupling the technifermions with the SM fermions either via heavy gauge bosons of an Extended Technicolor (ETC) sector [20–23] or through scalar fields in bosonic technicolor (BTC) [24–28]. In case the scalar mediators are elementary, one can control fine-tuning at scales above the elementary scalar masses by introducing supersymmetry (SUSY) [29,30]. Combining SUSY with TC is therefore appealing because of two general features

- The fundamental Higgs fields do not participate in electroweak symmetry breaking, but serve as natural messengers between the symmetry breaking sector and SM matter fields.
- The strong TC dynamics responsible for the EW symmetry breaking alleviates the little hierarchy problem of supersymmetric scenarios.

Recently supersymmetric TC models based on $\mathcal{N} = 4$ super Yang–Mills were considered in [31,32] where the low energy effective theory at scales below the TC scale, Λ_{TC} , was taken to be Minimal Walking Technicolor (MWT) [14,15,33]. In this paper we study a simpler but still, as we shall show, phenomenologically viable possibility within the UV complete theory of [31,32]: we assume a larger portion of the supersymmetric spectrum to be heavy which still allows the resulting low energy effective theory, featuring an $\text{SU}(3)$ global symmetry in the TC sector, to break correctly EW symmetry. Due to its global symmetry and the near conformal TC coupling we call this $\text{SU}(3)$ Walking Technicolor (3WT).

The paper is structured as follow: In Section 2 we review the particle content and renormalizable Lagrangian of MSCT, which represents the elementary description of 3WT. Next we integrate out the heavy mass eigenstates of MSCT and derive the effective Lagrangian at scales below the SUSY breaking scale, m_{SUSY} , in Section 3. In the subsequent Section 4 we write the effective Lagrangian at scales below Λ_{TC} , where the TC interaction becomes strong and is assumed to bind technifermions and technigluons within composite states. We continue in Sec-

Table 1

Non-MSSM superfield content of MSCT. Here Adj and \square denote the adjoint and fundamental representations, respectively. None of the fields above is charged under $SU(3)_c$.

Superfield	$SU(2)_{TC}$	$SU(2)_L$	$U(1)_Y$
Φ_L	Adj	\square	1/2
Φ_3	Adj	1	-1
V	Adj	1	0
Λ_L	1	\square	-3/2
N	1	1	1
E	1	1	2

tion 5 by working out the mass eigenstates, light Higgs couplings, and EW precision parameters of the model. Then, in Section 6, we test the 3WT viability by scanning the parameter space for data points satisfying the direct search limits on new particles and then performing a goodness of fit analysis of Higgs physics data at LHC with the viable scanned data points. The result of this analysis, which is that 3WT fits the current experimental data at a goodness level comparable to that of the SM, is the main result of the paper. Finally, we offer our conclusions in Section 7.

2. An UV complete technicolor model

The UV complete supersymmetric theory which provides our starting point is the same which has been introduced in [31,32] and called Minimal Supersymmetric Conformal Technicolor (MSCT). The gauge symmetry group of MSCT extends the SM one by the TC gauge group, $SU(2)_{TC}$. The TC vector superfields can be rearranged with the chiral superfields, containing the technifermions which transform under the adjoint representation of $SU(2)_{TC}$, in $\mathcal{N} = 4$ superfields. The MSCT Lagrangian can hence be expressed in compact form as that of the Minimal Supersymmetric Standard Model (MSSM) extended by an $\mathcal{N} = 4$ Super Yang Mills (4SYM) sector, which contains the TC sector. We furthermore add to the superpotential a fourth lepton superfamily sector, in which the fermion components are needed to cancel the Witten topological anomaly generated by the odd number of left-handed technifermions. The non-MSSM superfields introduced in MSCT and their quantum numbers are summarized in Table 1. Analogously the MSCT superpotential can be expressed in compact form in terms of the MSSM superpotential and its extension

$$P = P_{MSSM} + P_{TC}, \quad (1)$$

where P_{MSSM} is the MSSM superpotential, and P_{TC} is expressed by

$$P_{TC} = -\frac{g_{TC}}{\sqrt{2}} \epsilon^{abc} \Phi_L^a \cdot \Phi_L^b \Phi_3^c + y_U \Phi_L^a \cdot H_u \Phi_3^a + y_N \Lambda_L \cdot H_u N + y_E \Lambda_L \cdot H_d E + y_R E \Phi_3^a \Phi_3^a, \quad (2)$$

with H_u (H_d) denoting the $Y = +1/2$ ($-1/2$) Higgs superfield. The dot (\cdot) indicates a contraction between the $SU(2)_L$ doublets with the antisymmetric two-index Levi-Civita tensor ϵ .

To the potential obtained from Eq. (1) we add the soft SUSY breaking terms of the MSSM as well as those corresponding to P_{TC} , with the latter expressed by:

$$\begin{aligned}
\mathcal{L}_{soft}^{TC} = & - \left[a_{TC} \epsilon^{abc} \hat{U}_L^a \hat{D}_L^b \hat{U}_R^{*c} + a_U \hat{Q}_L^a \cdot \hat{H}_u \hat{U}_R^{*a} + a_N \hat{\Lambda}_L \cdot \hat{H}_u \hat{N}_R^* \right. \\
& + a_E \hat{\Lambda}_L \cdot \hat{H}_d \hat{E}_R^* + a_R \hat{U}_R^{*a} \hat{U}_R^{*a} \hat{E}_R^* + \frac{1}{2} M_D D_R^{\dagger a} D_R^{\dagger a} + c.c. \left. \right] - M_Q^2 \hat{Q}_L^{\dagger a} \hat{Q}_L^a \\
& - M_U^2 \hat{U}_R^{*a} \hat{U}_R^a - M_L^2 \hat{L}_L^\dagger \hat{L}_L - M_N^2 \hat{N}_R^* \hat{N}_R - M_E^2 \hat{E}_R^* \hat{E}_R,
\end{aligned} \tag{3}$$

where we write a hat on top of the scalar component of the chiral supermultiplets.

The model defined by Eqs. (1), (2) and (3) constitutes the fundamental description of the theory we study in this paper. The relevant scales of the model are the SUSY breaking scale, m_{SUSY} , and the EW scale that we identify with the low-energy strongly coupled regime of the TC theory $\Lambda_{TC} \sim 4\pi v_w$, which for $v_w = 246$ GeV implies $\Lambda_{TC} \sim 3$ TeV. We will assume here that the two scales satisfy

$$m_{SUSY} \gg \Lambda_{TC}. \tag{4}$$

With this ordering the EW symmetry is broken dynamically. Furthermore, we assume the mass spectrum to feature roughly the following hierarchy:

- 1) All SUSY breaking masses as well as the μ parameter are of $O(m_{SUSY})$, therefore all the superpartners as well as the elementary Higgs scalars have masses of the same order.
- 2) All the lightest composite states acquire masses, which are at most of the order of Λ_{TC} .

In the following section we proceed to derive the effective Lagrangian describing the physics below the SUSY breaking scale by integrating out the heavy states.

3. Mesoscopic Lagrangian

After we integrate out all particles with mass greater than m_{SUSY} , the only technifermions left are the fermionic components of the EW doublet superfield Φ_L and EW singlet Φ_3 in Table 1:

$$Q_L^a = \begin{pmatrix} U_L^a \\ D_L^a \end{pmatrix}, \quad U_R^a, \quad a = 1, 2, 3. \tag{5}$$

To derive the effective Lagrangian, valid between the scales Λ_{TC} and m_{SUSY} , one first writes down the Higgs Yukawa sector in MSCT:

$$\begin{aligned}
-\mathcal{L}_{Yukawa}^{MSCT} = & \hat{H}_u \cdot F_u + \hat{H}_d \cdot F_d + \text{h.c.}, \\
F_u = & q_{Lu}^i Y_u^i u_R^{\dagger i} + y_U Q_L U_R^\dagger + y_N L_L N_R^\dagger, \\
F_d = & q_{Ld}^i Y_d^i d_R^{\dagger i} + l_L^i Y_l^i e_R^{\dagger i} + y_E L_L E_R^\dagger,
\end{aligned} \tag{6}$$

where $i = 1, 2, 3$ is the flavor index and it is summed over. The matrices Y_u , Y_d , and Y_l are diagonal, and the CKM matrix V is contained in the definitions of the vectors

$$q_{Lu}^{Ti} = (u_L^i, V^{ij} d_L^j) \quad \text{and} \quad q_{Ld}^{Ti} = (V^{\dagger ij} u_L^j, d_L^i). \tag{7}$$

Given that the potential of the MSSM Higgs fields is

$$V_{MSSM} = \left(m_{SUSY}^2 + |\mu|^2 \right) |\hat{H}_u|^2 + \left(m_{SUSY}^2 + |\mu|^2 \right) |\hat{H}_d|^2 - \left(b \hat{H}_u \hat{H}_d + \text{h.c.} \right) + \dots \tag{8}$$

by solving the equation of motion in terms of the Higgs mass eigenstates and plugging the solutions back into Eq. (6), leads us to the first line of the following dimension six interaction terms for the fermions in the intermediate scale (or mesoscopic) effective Lagrangian

$$\begin{aligned} \mathcal{L}_{4\text{-fermion}} = & \frac{c_\theta^2}{m_s^2} \left(F_u^\dagger F_u + F_d^\dagger F_d \right) - \frac{c_\theta s_\theta}{m_s^2} (F_u \cdot F_d + \text{h.c.}) \\ & + \frac{g_{\text{TC}}^2}{m_{\text{SUSY}}^2} \epsilon_{abc} \epsilon_{cde} \eta_i^{\alpha a} \eta_{j\alpha}^b \eta_{i\beta}^{\dagger d} \eta_j^{\dagger \beta e} . \end{aligned} \quad (9)$$

Similarly, the last term originates in an analogous way from the potential and the Yukawa sector of 4SYM. In Eq. (9) we have defined

$$\eta_\alpha^T = \left(U_{L\alpha}, D_{L\alpha}, -i\sigma_{\alpha\beta}^2 U_R^{\dagger\beta} \right) , \quad (10)$$

where σ^2 is the second Pauli matrix, and the indices i and j denote SU(3) flavor; the first letters of the alphabet are reserved for the adjoint SU(2) technicolor indices, while the Greek indices label the spin component, and the TC indices, running from 1 to 3, are written explicitly only in the last term. We suppress summed spin indices as long as it can be done consistently. Finally, we have defined

$$m_s^2 = \left(\mu^2 + m_{\text{SUSY}}^2 \right) \frac{(\mu^2 + m_{\text{SUSY}}^2)^2 - b^2}{(\mu^2 + m_{\text{SUSY}}^2)^2 + b^2} , \quad \tan \theta = \frac{b}{\mu^2 + m_{\text{SUSY}}^2} . \quad (11)$$

In the rest of this paper we use abbreviations $s_\theta \equiv \sin \theta$, $c_\theta \equiv \cos \theta$ and $t_\theta \equiv \tan \theta$.

The four-fermion interaction terms in Eq. (9) are relevant because they eventually give mass to the SM fermions once the technifermions condense. The first line in Eq. (9) derives from decoupling the Higgs scalars, and breaks the global SU(3) symmetry, while the last term in Eq. (9), stemming from the 4SYM sector, respects the global SU(3) symmetry of the pure TC sector.

At energy scales below Λ_{TC} the TC interaction becomes strong, and physical states charged under TC get bound in composite states with zero TC charge. In the next section therefore we derive the effective Lagrangian involving such states.

4. Effective Lagrangian at the electroweak scale

Similarly to QCD, in TC a tower of composite states is predicted to arise at low energies. At scales below Λ_{TC} the new physics degrees of freedom are the composite states associated with the strong TC interaction, and the form of the effective Lagrangian is constrained to satisfy the approximate global symmetries of the fundamental Lagrangian. In the following we derive the effective Lagrangian introducing first the composite scalars and then the composite vectors.

4.1. Technicolor scalar sector

The composite scalar matrix field M , singlet under $\text{SU}(2)_{\text{TC}}$, has minimal particle content given by the techniquark bilinears:

$$M_{ij} \sim \mu_i^\alpha \eta_j^\beta \varepsilon_{\alpha\beta} = \eta_i \eta_j , \quad \text{with} \quad i, j = 1 \dots 3. \quad (12)$$

The field M transforms under the full SU(3) group according to

$$M \rightarrow u M u^T , \quad \text{with} \quad u \in \text{SU}(3) . \quad (13)$$

The effective linearly transforming SU(3) invariant Lagrangian reads:

Table 2

Transformation properties of the component fields of the matrix M under $SU(2)_L \times U(1)_Y$. The complex scalars are grouped, based on their transformation properties under $SU(2)_L$, into one triplet, one doublet, and one singlet.

Field	$SU(2)_L$	$U(1)_Y$
$\Delta \sim Q_L Q_L$	$\square\square$	1
$\sigma \sim Q_L U_R^\dagger$	\square	$-\frac{1}{2}$
$\delta^{--} \sim U_R^\dagger U_R^\dagger$	1	-2

$$\mathcal{L}_M = \frac{1}{2} \text{Tr} \left[D_\mu M^\dagger D^\mu M \right] - \mathcal{V}_M, \quad (14)$$

where the covariant derivative is given by

$$D_\mu M = \partial_\mu M - i g_L \left[G_\mu M + M G_\mu^T \right],$$

with

$$G_\mu = \tilde{W}_\mu^a \frac{\lambda^a}{2} + t_\xi B_\mu Y_M, \quad a = 1, 2, 3. \quad (15)$$

In the above equation λ^a are the Gell-Mann matrices, $t_\xi = \tan \xi$ with ξ the EW mixing angle, \tilde{W}_μ and B_μ are the SM EW gauge fields, and

$$Y_M = \text{diag} \left(\frac{1}{2}, \frac{1}{2}, -1 \right). \quad (16)$$

The most general $SU(3)$ preserving effective potential, including operators up to dimension four, is²

$$\mathcal{V}_M = -\frac{m^2}{2} \text{Tr} \left[M^\dagger M \right] + \frac{\lambda}{4} \text{Tr} \left[M^\dagger M \right]^2 + \lambda' \text{Tr} \left[M^\dagger M M^\dagger M \right] - 2m' \left[\det M + \det M^\dagger \right], \quad (17)$$

which breaks $SU(3)$ spontaneously to $SO(3)$ for positive m^2 , as we show explicitly in [Appendix A](#). The TC gauge interaction is actually invariant under $U(3) \equiv SU(3) \times U(1)_A$, rather than just $SU(3)$. However the $U(1)_A$ axial symmetry is anomalous, and is therefore broken at the quantum level. This corresponds to the $\det M$ term in Eq. (17). The components of the matrix $M \sim \eta^T \eta$ can be described in terms of the transformation properties of the composite states under $SU(2)_L \times U(1)_Y$. This notation is introduced in [Table 2](#).

Using this notation, the matrix M is written in terms of complex scalars as

$$M = \begin{pmatrix} \sqrt{2}\Delta^{++} & \Delta^+ & \sigma^0 \\ \Delta^+ & \sqrt{2}\Delta^0 & \sigma^- \\ \sigma^0 & \sigma^- & \sqrt{2}\delta^{--} \end{pmatrix}. \quad (18)$$

This notation is suitable to study the vacuum, since the flavor extension sector breaks the global symmetry of the potential from $SU(3)$ down to the EW gauge group $SU(2) \times U(1)$. Next, we discuss how to consistently introduce also the composite vector fields.

² In principle the higher dimensional operators can play a role and should be systematically included. For an initial investigation and qualitative account of the various constraints, we truncate the effective theory at the level of dimension four operators. This provides a quantitative baseline for possibly more refined analyses in the future.

4.2. Vector sector

A minimal set of composite vector fields transforming homogeneously under SU(3) can be written, in terms of Gell-Mann matrices λ^a , as

$$A_\mu = A_\mu^a \frac{\lambda^a}{2}, \quad (19)$$

which transform under SU(3) according to

$$A_\mu \rightarrow u A_\mu u^\dagger, \quad \text{with} \quad u \in \text{SU}(3). \quad (20)$$

The elementary particle content of A^μ is expressed by the equivalence

$$\begin{aligned} A_i^{\mu j} &\sim \eta_i^\alpha \sigma_{\alpha\beta}^\mu \eta^{\dagger\beta j} - \frac{1}{3} \delta_i^j \eta_k^\alpha \sigma_{\alpha\beta}^\mu \eta^{\dagger\beta k} = \bar{\Psi}^j \gamma^\mu \Psi_i - \frac{1}{3} \delta_i^j \bar{\Psi}^k \gamma^\mu \Psi_k, \\ \bar{\Psi} &= (\bar{U}_L, \bar{D}_L, \bar{U}_R^c), \end{aligned} \quad (21)$$

where the components of Ψ in SU(3) space are Dirac spinors, with the superscript c on the last entry denoting the charge conjugation. The vector and axial-vector charge eigenstates and their elementary particle content are given in [Appendix B](#).

The effective Lagrangian including composite vector fields, A^μ , has already been derived in [\[33\]](#) for a theory with SU(4) global symmetry in the TC sector by applying the hidden local symmetry principle [\[34,35\]](#). Those results can be straightforwardly used for SU(3) symmetric TC by defining the corresponding vector field and the field strength tensor:

$$C_\mu = A_\mu - \epsilon G_\mu, \quad \epsilon = \frac{gL}{g_{\text{TC}}}, \quad F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu - i g_{\text{TC}} [A_\mu, A_\nu], \quad (22)$$

with G_μ defined in Eq. [\(15\)](#). The vector field C_μ has the same transformation law as A_μ :

$$C_\mu \rightarrow u C_\mu u^\dagger, \quad \text{with} \quad u \in \text{SU}(3). \quad (23)$$

The kinetic and mass terms for the vector fields can then be written as

$$\mathcal{L}_V = -\frac{1}{2} \text{Tr} [\tilde{W}_{\mu\nu} \tilde{W}^{\mu\nu}] - \frac{1}{4} B_{\mu\nu} B^{\mu\nu} - \frac{1}{2} \text{Tr} [F_{\mu\nu} F^{\mu\nu}] + m_A^2 \text{Tr} [C_\mu C^\mu], \quad (24)$$

while the scalar-vector field interaction terms up to dimension four operators read

$$\begin{aligned} \mathcal{L}_{M-V} &= g_{\text{TC}}^2 r_1 \text{Tr} [C_\mu C^\mu M M^\dagger] + g_{\text{TC}}^2 r_2 \text{Tr} [C_\mu M C^{\mu T} M^\dagger] \\ &\quad - g_{\text{TC}}^2 \frac{r_3}{4} \text{Tr} [C_\mu C^\mu] \text{Tr} [M M^\dagger], \end{aligned} \quad (25)$$

with constants $r_i \sim \mathcal{O}(1)$. A few remarks are in order: First, higher dimensional operators are suppressed by powers of Λ_{TC} , and are therefore subleading. Second, terms proportional to y_U , which explicitly break SU(3) global symmetry, are small compared to those proportional to g_{TC}^2 and can therefore be neglected at leading order. Third, to simplify the phenomenological analysis of 3WT, presented in the next section, we neglect also a covariant derivative coupling term (see [\[33\]](#) for its precise definition).³ Finally, in the next subsection, we determine the effective Lagrangian terms of the flavor extension of 3WT below scale Λ_{TC} and then summarize the 3WT full Lagrangian.

³ Neglecting this term is simply a restriction on the parameter space: this term could be included in more thorough future analyses.

4.3. Flavor extension sector

The four-fermion theory, Eq. (9), is given just below the SUSY breaking scale and the techniquark condensate needs to be evolved down to the EW scale. This is achieved by multiplying the techniquark Yukawa coupling y_U , renormalized at the SUSY breaking scale, with the dimensionless factor

$$\omega = \frac{\langle U_L U_R^\dagger \rangle_{m_{\text{SUSY}}}}{\langle U_L U_R^\dagger \rangle_{\Lambda_{\text{TC}}}} = \left(\frac{m_{\text{SUSY}}}{\Lambda_{\text{TC}}} \right)^\gamma, \quad (26)$$

written under the assumption that the anomalous dimension γ of the techniquark mass operator is constant.

Note that in the following we neglect the contribution of the last term in Eq. (9) because that term respects the global SU(3) symmetry, and therefore its effects should already be parametrized by the quartic couplings in the TC effective Lagrangian, Eq. (17). The masses of the SM fermions and the fourth family leptons arise from the terms on the first line of Eq. (9): more specifically those masses are generated by the following four-fermion operator

$$\eta^T K \eta, \quad (27)$$

with

$$\begin{aligned} K_{ij} = & \frac{y_U c_\theta \omega}{m_s^2} \left[\delta_{ik} c_\theta \left(q_{Lu}^{\dagger k} Y_u^* u_R + y_N^* L_L^{\dagger k} N_R \right) \right. \\ & \left. - \epsilon_{ik} s_\theta \left(q_{Ld}^k Y_d d_R^\dagger + l_L^k Y_l e_R^\dagger + y_E L_L^k E_R^\dagger \right) \right] \delta_{3j}, \\ & i, j = 1, \dots, 3; \quad k = 1, 2; \quad \epsilon_{3k} \equiv 0, \end{aligned} \quad (28)$$

upon condensation of the techniquarks. Under SU(3) global symmetry the spurion K transforms as $K \rightarrow u^* K u^\dagger$.

The four-techniquark term on the other hand is

$$\begin{aligned} \frac{y_U^2 c_\theta^2}{m_s^2} \omega^2 (Q_L U_R^\dagger) (Q_L^\dagger U_R) &= K'_{ijkl} \eta_i^\alpha \eta_{j\alpha} \eta_{k\beta}^\dagger \eta_l^{\dagger\beta}, \\ K'_{ijkl} &= \frac{y_U^2 c_\theta^2}{m_s^2} \omega^2 (\delta_{ik1} + \delta_{ik2}) \delta_{jl3}, \end{aligned} \quad (29)$$

where α and β are spin indices. For this term to be invariant under SU(3), the spurion K' must transform as $K'_{ijkl} \rightarrow u_{im} u_{jn} u_{ko}^* u_{lp}^* K'_{mnop}$, with $u \in \text{SU}(3)$. To estimate the effects of renormalization, we simply assume factorization, leading to a multiplicative factor of ω^2 .

At the lowest order in the spurions, the SU(3) breaking effective Lagrangian, obtained from Eqs. (28) and (29) is:

$$\mathcal{L}_F = c_1 \Lambda_{\text{TC}}^2 \text{Tr}[MK] + c_2 \Lambda_{\text{TC}}^4 K'_{ijkl} M_{ij} M_{kl}^* + \text{h.c.}, \quad (30)$$

where we introduced factors of Λ_{TC} to define the dimensionless coefficients c_i , which parametrize the couplings of the effective Lagrangian in terms of those of the underlying theory. We estimate these coefficients using dimensional analysis [36–38] and find

$$c_1 = \mathcal{O}(\Upsilon^{-1}), \quad c_2 = \mathcal{O}(\Upsilon^{-2}), \quad \Upsilon \equiv \frac{\Lambda_{\text{TC}}}{v_w}. \quad (31)$$

To summarize, the 3WT effective Lagrangian is given by

$$\mathcal{L}_{3WT} = \mathcal{L}_k + m_A^2 \text{Tr}[C_\mu C^\mu] + \mathcal{L}_{M-V} + \mathcal{L}_{\text{Yukawa}} - \mathcal{V}, \quad (32)$$

where \mathcal{L}_k contains all the covariant kinetic terms for the composite scalar and vector fields given respectively in Eqs. (14) and (24). The terms included in \mathcal{L}_{M-V} are given by Eq. (25), while the Yukawa interactions and the potential are defined by

$$\mathcal{L}_{\text{Yukawa}} = c_1 \Lambda_{\text{TC}}^2 \text{Tr}[MK] + \text{h.c.}, \quad \mathcal{V} = \mathcal{V}_M - \left[c_2 \Lambda_{\text{TC}}^4 K'_{ijkl} M_{ij} M_{kl}^* + \text{h.c.} \right], \quad (33)$$

with TC potential \mathcal{V}_M given by Eq. (17). Having at hand the full Lagrangian for the model, we proceed to work out its phenomenological consequences in the next section.

5. Phenomenology

In this section we work out the mass eigenstates arising at low energy in 3WT, the couplings of the light composite Higgs boson to SM mass eigenstates, and the EW precision parameters S and T .

5.1. EW symmetry breaking & mass eigenstates

The most general ground state which breaks the EW gauge group down to the electromagnetic $U(1)_Q$ can be parametrized by the following form of the vacuum expectation value (vev) of the matrix field M

$$\langle M \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & 0 & v_\sigma \\ 0 & \sqrt{2}v_\Delta & 0 \\ v_\sigma & 0 & 0 \end{pmatrix}. \quad (34)$$

Minimizing the scalar potential given in Eqs. (17) and (30) leads to the following relations between the parameters:

$$m^2 = \lambda (v_\sigma^2 + v_\Delta^2) + 4\lambda' v_\Delta^2 - 2\lambda'' v_\sigma^2, \quad \tilde{m}^2 = 2(\lambda' + \lambda'') (v_\sigma^2 - 2v_\Delta^2), \quad (35)$$

where

$$\tilde{m}^2 = c_2 y_U^2 c_\beta^2 \frac{\Lambda_{\text{TC}}^4 \omega^2}{m_s^2}, \quad \lambda'' \equiv -\frac{m'}{v_\Delta}. \quad (36)$$

A sufficient condition for the potential to be bounded from below is $\lambda, \lambda' > 0$. In the limit $v_\sigma = \sqrt{2}v_\Delta$ we obtain $\tilde{m} = 0$, consistently with the pure TC result in Appendix A. We notice that the ground state changes because of the four-technifermion interaction in Eq. (29). An analogous change of ground state occurs in ETC theories where some of the chiral symmetries of the pure TC theory are broken by extended gauge interactions. In our model setup the effective four fermion interactions at low energy arise from attractive Yukawa couplings which is different from the usual ETC scenario where the underlying gauge interactions can be either repulsive or attractive.

The mass spectrum of 3WT, corresponding to the vev in Eq. (34), includes two neutral scalars, h^0 and H^0 , as well as one neutral pseudoscalar, Π^0 , one singly charged and two doubly charged scalars, H^\pm , $h^{\pm\pm}$, and $H^{\pm\pm}$, respectively, with the corresponding squared mass matrices given in Appendix C. We introduce for later use the mixing angle φ defined by

$$\begin{pmatrix} h^0 \\ H^0 \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} c_\varphi & -s_\varphi \\ s_\varphi & c_\varphi \end{pmatrix} \begin{pmatrix} \Re(\sigma^0) \\ \Re(\Delta^0) \end{pmatrix}. \quad (37)$$

Assuming that only the third generation SM fermions have non-negligible Yukawa couplings, it follows from Eq. (30) that the masses of the upper component u and the lower component d of a generic SM fermion EW doublet are given, respectively, by

$$m_u = c_1 \frac{c_\theta^2 y_U y_u \omega \Lambda_{\text{TC}}^2}{m_s^2} \frac{v_\sigma}{\sqrt{2}}, \quad m_d = \frac{y_d}{y_u} t_\theta m_u. \quad (38)$$

Using the previous equation, the fact that both members of the fourth fermion doublet, E and N , have to be heavier than about 100 GeV, and requiring the Yukawa couplings to be perturbative, we deduce that θ cannot be close to either 0 or $\pi/2$. To simplify the study of 3WT from here on we take $m_s = m_{\text{SUSY}}$, which from Eq. (11) is equivalent to imposing:

$$m_{\text{SUSY}} = m_s \quad \Rightarrow \quad \frac{b}{\mu^2} = \frac{t_\theta + t_\theta^{-1}}{2}. \quad (39)$$

From the fact that neither t_θ nor its inverse are large, it follows that our choice does not introduce any large hierarchy between b and μ^2 , and both can be taken of $\mathcal{O}(m_s^2)$. This is consistent with our assumption that the higgsinos are heavy and decoupled from low energy phenomenology.

Finally, the composite axial-vector and vector resonances mix with the SM gauge bosons, while the doubly charged baryon technivector does not. The resulting physical massive vector states are $Z_\mu, Z'_\mu, Z''_\mu, W_\mu^\pm, W'_\mu^\pm, W''_\mu^\pm, \Omega_\mu^{\pm\pm}$, with the corresponding squared mass matrices given in Appendix C. The masses of W^\pm and Z in the limit of negligible mixing ($\epsilon = 0$) read:

$$m_W^2 \cong \frac{1}{4} g_L^2 v_w^2, \quad m_Z^2 \cong \frac{1}{4} (g_L^2 + g_Y^2) (1 + t_\rho^2) v_w^2, \quad (40)$$

where $t_\rho = \sqrt{2} v_\Delta / v_\sigma$, and the EW scale is given by

$$v_w^2 = (\sqrt{2} G_F)^{-1} = (246 \text{ GeV})^2 = v_\sigma^2 + 2v_\Delta^2. \quad (41)$$

In the next subsection we give a more detailed discussion on the Higgs mass.

5.2. Higgs mass

The mass of a light Higgs in TC models deserves a little introduction. In [17] it has been shown that $m_{h^0}^2$ receives a large negative corrections at one loop from the top quark and the W^\pm and Z bosons. In [18,19] these considerations have been extended to fully dynamical settings of electroweak symmetry breaking and generation of fermion masses. Because of these corrections the tree level Higgs mass can be significantly larger than 126 GeV: adapting the formulas given in [17] to the 3WT case, while neglecting in first approximation any mixing between the two neutral scalar states, we can write

$$\left(m_{h^0}^2\right)_{\text{tree}} \simeq m_{h^0}^2 + \frac{8}{3} \kappa^2 \left[2a_f^2 (3m_t^2 + m_E^2 + m_N^2) - 3a_\pi^2 \left(m_W^2 + \frac{m_Z^2}{2} \right) \right], \quad (42)$$

where $(m_{h^0}^2)_{\text{tree}}$ is the scalar mass due to pure strong dynamics, neglecting the scalar coupling with the EW gauge currents and SM matter. Furthermore, κ is an $\mathcal{O}(1)$ renormalization coefficient, while a_f and a_π are rescaling coefficients of the SM Higgs linear coupling to fermions and quadratic coupling to gauge bosons, respectively. In writing the formula above we set the

dimension of the TC representation to 3 and the number of technidoublets to 0.75, as it is the case for 3WT. To illustrate the idea we evaluate the formula above by assuming couplings and renormalization coefficient to be SM like ($a_F = a_\pi = \kappa = 1$), and $m_E = m_N = m_t$. With these choices we find a light Higgs, $(m_{h^0}^2)_{\text{tree}} = 740 \text{ GeV}$. This value is not far from the naive TC estimate for 3WT, obtained by scaling up the mass of the $f_0(500)$ QCD resonance [39] ($m_{f_0} = 400\text{--}550 \text{ GeV}$):

$$(m_{h^0}^2)_{\text{naive}} \simeq \frac{4}{3} \frac{v_w^2}{f_\pi^2} m_{f_0}^2 = 1200\text{--}1700 \text{ GeV} . \quad (43)$$

The required suppression of the tree level Higgs mass, around 50%, might realistically come from near-conformal dynamics. This should be indeed the case for 3WT, which has 1.5 adjoint Dirac flavors and is outside but close to the conformal window [40].

As this discussion implies, there will be some amount of fine tuning involved in this scenario. However, the same applies to practically any BSM scenario which explains the relatively small mass of the observed Higgs boson. In our case the amount of fine tuning can be quantified in a simple way as

$$\text{FT} = \frac{(125 \text{ GeV})^2}{(m_{h^0}^2)_{\text{naive}}} . \quad (44)$$

The concrete values are on the level of few percents. For example, for $(m_{h^0})_{\text{naive}} \simeq 1 \text{ TeV}$, $\text{FT} \simeq 1.6\%$, whereas for $(m_{h^0})_{\text{naive}} \sim 740 \text{ GeV}$, $\text{FT} \simeq 2.9\%$.

We observe that the couplings involved in the one loop correction to m_{h^0} are SM-like in size, and this ensures that higher order corrections are indeed perturbative and do not spoil the partial cancellation with the effective TC tree level contribution.

5.3. Coupling coefficients

In our model the linear Higgs coupling to charged vector bosons can be written in compact form as

$$\mathcal{L} \supset \frac{2m_A^2}{v_w} \bar{W}_\mu^\dagger \Xi \bar{W}^\mu h^0 , \quad \bar{W}_\mu^\dagger = \left(\tilde{W}_\mu^-, V_\mu^-, A_\mu^-, \Omega_\mu^{--} \right) , \quad (45)$$

with the non-zero terms of the matrix Ξ given by

$$\begin{aligned} \Xi_{1,1} &= \left(x^2 + \epsilon^2 z_1 \right) \left(c_\varphi c_\rho - \sqrt{2} s_\varphi s_\rho \right) + \frac{\epsilon^2 z_3}{\sqrt{2}} \left(s_\varphi s_\rho - \sqrt{2} c_\varphi c_\rho \right) , \\ \Xi_{2,2} &= z_1 \left(c_\varphi c_\rho - \sqrt{2} s_\varphi s_\rho \right) + z_2 \left(c_\varphi s_\rho - \sqrt{2} c_\rho s_\varphi \right) + \frac{z_3}{\sqrt{2}} \left(s_\varphi s_\rho - \sqrt{2} c_\varphi c_\rho \right) , \\ \Xi_{3,3} &= z_1 \left(c_\varphi c_\rho - \sqrt{2} s_\varphi s_\rho \right) - z_2 \left(c_\varphi s_\rho - \sqrt{2} c_\rho s_\varphi \right) + \frac{z_3}{\sqrt{2}} \left(s_\varphi s_\rho - \sqrt{2} c_\varphi c_\rho \right) , \\ \Xi_{1,2} &= -\frac{\epsilon}{2\sqrt{2}} \Xi_{2,2} , \quad \Xi_{1,3} = -\frac{\epsilon}{2\sqrt{2}} \Xi_{3,3} , \\ \Xi_{4,4} &= 2c_\varphi c_\rho \left(z_1 + z_2 \right) + \frac{z_3}{\sqrt{2}} \left(s_\varphi s_\rho - \sqrt{2} c_\varphi c_\rho \right) . \end{aligned} \quad (46)$$

Here φ is the mixing angle of the neutral scalars, Eq. (37), and

$$\epsilon = \frac{gL}{g_{\text{TC}}}, \quad x = \frac{gLv_w}{2m_A}, \quad t_\rho = \frac{\sqrt{2}v_\Delta}{v_\sigma}, \quad z_i = \left(\frac{g_{\text{TC}}v_w}{2m_A} \right)^2 r_i, \quad i = 1, 2, 3. \quad (47)$$

To simplify the analysis we select the slice of parameter space where the axial vector coupling to the light Higgs is zero ($\Xi_{1,3} = \Xi_{3,3} = 0$). Moreover, we require the mixing mass term of the axial-vector to arise only from mixing ($(\mathcal{M}_W^2)_{1,3} = -\epsilon m_A^2/\sqrt{2}$), which is taken to be small. Together these conditions are satisfied by imposing

$$z_1 = \frac{1}{4} \left(3 + c_{2\rho}^{-1} \right) z_3, \quad z_2 = \frac{1}{4} t_{2\rho} z_3. \quad (48)$$

With the help of the two equations above, the charged vector boson squared mass matrix in Eq. (C.6) can be written in compact form as [41]:

$$\mathcal{M}_W^2 = \begin{pmatrix} m_{\tilde{W}}^2 & -\frac{\epsilon m_V^2}{\sqrt{2}} & -\frac{\epsilon m_A^2}{\sqrt{2}} \\ -\frac{\epsilon m_V^2}{\sqrt{2}} & m_V^2 & 0 \\ -\frac{\epsilon m_A^2}{\sqrt{2}} & 0 & m_A^2 \end{pmatrix}, \quad (49)$$

with

$$m_{\tilde{W}}^2 = \left[x^2 + (1 + s^2) \epsilon^2 \right] m_A^2, \quad m_V^2 = (1 + 2s^2) m_A^2, \quad s^2 = \frac{z_3}{4} s_{2\rho} t_{2\rho}. \quad (50)$$

The mass eigenvalues can be expanded in x and ϵ , which in TC are both expected to be small:

$$\begin{aligned} m_{\tilde{W}}^2 &\cong m_A^2 x^2 \left[1 - \epsilon^2 \right], & m_{W''}^2 &\cong m_A^2 \left[1 + \frac{1}{2} (1 + x^2) \epsilon^2 - \frac{1}{8} \left(2 + \frac{1}{s^2} \right) \epsilon^4 \right], \\ m_{W'}^2 &\cong m_A^2 \left[1 + 2s^2 + \frac{1}{2} (1 + 2s^2 + x^2) \epsilon^2 + \frac{1}{8} \left(2 + \frac{1}{s^2} \right) \epsilon^4 \right], \end{aligned} \quad (51)$$

where we kept contributions up to $\mathcal{O}(x^n \epsilon^{4-n})$, with $n = 0, \dots, 4$.

Generally the linear couplings of the light Higgs are conveniently expressed in terms of coupling coefficients defined by

$$\begin{aligned} \mathcal{L}_{\text{eff}} &= \sum_i a_{W_i} \frac{2m_{W_i}^2}{v_w} h W_{i\mu}^+ W_i^{-\mu} + \sum_j a_{Z_j} \frac{m_{Z_j}^2}{v_w} h Z_{j\mu} Z_j^\mu \\ &\quad - a_f \sum_{\psi=t,b,\tau,N,E} \frac{m_\psi}{v_w} h \bar{\psi} \psi - \sum_k a_{S_k} \frac{2m_{S_k}^2}{v_w} h S_k^+ S_k^-, \end{aligned} \quad (52)$$

where the indices i and k run over all the charged scalars and vector bosons (including the states with double charge). By normalizing the linear Higgs couplings to the charged vector bosons according to Eq. (52), with the masses given by Eqs. (51), (C.7), we determine the relations

$$\begin{aligned} a_W &= \left(c_\varphi c_\rho - \sqrt{2} s_\varphi s_\rho \right) \left[1 - \frac{x^2 \epsilon^2}{2} \left(1 + \frac{1}{1 + 2s^2} \right) \right] + \frac{x^2 \epsilon^2 s^2}{2(1 + 2s^2)^2} \left(\frac{c_\varphi}{c_\rho} - \sqrt{2} \frac{s_\varphi}{s_\rho} \right), \\ a_{W'} + a_{W''} &= c_\varphi c_\rho - \sqrt{2} s_\varphi s_\rho + \frac{s^2}{1 + 2s^2} \left(\frac{c_\varphi}{c_\rho} - \sqrt{2} \frac{s_\varphi}{s_\rho} \right), \\ a_\Omega &= \frac{s^2}{s_\rho} \frac{2c_\varphi (1 + t_\rho^{-1}) + \sqrt{2} c_{2\rho} c_\rho^{-2} s_\varphi}{2 - 4s^2 + s^2 \left(2(t_\rho^{-1} + s_\rho^{-2}) + c_\rho^{-2} \right)}. \end{aligned} \quad (53)$$

The corresponding result for fermions reads simply

$$a_f = \frac{c_\varphi}{c_\rho} . \quad (54)$$

Among the neutral vector resonances only Z is relevant for the LHC observables we include in this study. As we discuss in the next section, the numerical values of a_Z that we find are very close to a_W , and for all practical purposes they can therefore be taken equal to each other.⁴

5.4. Oblique corrections

The precision EW parameters [42,43] can be calculated directly from the vector-boson sector of the effective Lagrangian, Eqs. (24) and (25), by integrating out the heavy charged and neutral states and then using the formulas provided by [44,45].⁵ At tree level and linear order in the mixing parameter ϵ we find:

$$S_{tree} = 0, \quad \alpha_e T_{tree} = -\frac{2v_\Delta^2}{v_w^2} . \quad (55)$$

For the T parameter to be consistent with the experiments, the vev component v_Δ clearly has to be small in comparison to the EW scale. We note that the S parameter is zero up to corrections of order ϵ^4 while the T parameter obtains further contributions of order ϵ^2 which we neglect since $\epsilon \ll 1$.

The intrinsic TC contribution is usually calculated from the one loop perturbative diagrams of the technifermions, which are assumed to be massive because of dynamical symmetry breaking. The dynamical mass divided by m_Z is usually taken infinite, as this gives a meaningful result with no unknown parameters. These contributions are denoted by S_{naive} and T_{naive} . For our underlying technicolor theory, Eq. (5), the dynamical masses are such that while the up-techniquark, U , gains only a Dirac mass, m_U , the down-techniquark, D_L , acquires also a Majorana mass, m_L . Oblique corrections for this general case have been calculated in [46] in terms of integral functions, which we use to derive the explicit formulas given in Appendix D. In the pure 3WT limit the Dirac and Majorana masses are equal to each other,⁶ $m_L = m_U$, and we find

$$S_{naive} \approx \frac{3}{4\pi}, \quad T_{naive} \approx \frac{m_U^2 \log[\frac{m_U}{\Lambda_{NP}}]}{4m_Z^2 s_\xi^2 c_\xi^2 \pi} . \quad (56)$$

The dependence of T on the renormalization scale Λ_{NP} is due to the Majorana mass, and it should be matched onto a renormalizable term in the underlying theory. Indeed in the full theory, Eqs. (1), (2), (3), the contribution to T of the technifermion D_R cancels the one in Eq. (56) when its mass is the same as that of U . Phenomenological viability therefore would require the two masses to be close in value, and we implement this requirement in the current analysis by assuming $m_D \sim \Lambda_{NP}$ and Λ_{NP} to be close to m_U . With this assumption the one-loop contribution T_{naive} is negligible when compared to the tree-level contribution, T_{tree} , and therefore we approximate the naive parameters as given by:

$$S_{naive} \approx \frac{3}{4\pi}, \quad T_{naive} \approx 0 . \quad (57)$$

⁴ The analytic expressions for a_{Z_j} and a_{S_k} are lengthy and we do not reproduce them here.

⁵ For this task we adapted the code provided by the authors of [33] to 3WT.

⁶ This follows from the SO(3) invariance of the mass terms implied by Eq. (18).

Another independent contribution to S and T in 3WT comes from the fourth family leptons N and E . These have been already evaluated in [47]: for $m_E > m_N$, the $S_{N,E}$ contribution is actually negative and can offset the positive S_{naive} . We can therefore summarize the non-negligible contributions to the S and T parameters, with the assumptions described above, as:

$$S = S_{naive} + S_{N,E} , \quad T = T_{tree} + T_{N,E} , \quad (58)$$

where T_{tree} , S_{naive} , and the Higgs contributions are given in Eqs. (55) and (57).

The contribution of the heavy leptons to the S parameter is indeed critical to partially cancel the naive TC contribution (~ 0.24) and match the experimental result (0.04 ± 0.09): the required cancellation imposes a moderate constraint on the ratio m_E/m_N [15,46].

In the next section we study the viability of 3WT by performing a goodness of fit analysis based on the recent LHC and Tevatron data on Higgs physics, as well as the experimental values of the S and T parameters.

6. Experimental validation

To test the model's viability, we start by performing a numerical scan over the parameter space and collecting data points that satisfy the direct search lower bounds on the new physics mass spectrum and the EW precisions tests. We also require perturbativity of all couplings excluding g_{TC} , and stability of the potential at large values of the scalar fields. The unknown parameters derived from strong dynamics are estimated using dimensional analysis [36–38]. Finally, we fix the anomalous dimension to $\gamma = 1.5$, and require the supersymmetry breaking scale to be larger than 5 TeV. Putting all this together, the free parameters in our scan acquire values in the domain defined by the following relations:

$$\begin{aligned} 210 \text{ GeV} \leq |v_\sigma| \leq 246 \text{ GeV} , \quad m_A = 1 \text{ TeV} , \quad \gamma = 1.5 , \quad \pi \leq \Upsilon \leq 4\pi , \quad 0.5 \leq c_1 \Upsilon^{-1} \leq 5 , \\ 0.1 \leq \lambda \leq (2\pi)^2 , \quad 0.1 \leq y_t, y_N, y_E, y_U \leq 2\pi , \quad 0.1 \leq \lambda'' \leq 200 , \quad 0 \leq z_3 \leq 1 , \quad |\epsilon| \leq 2x . \end{aligned} \quad (59)$$

The remaining model parameters, λ' , θ , x and v_Δ , are determined in terms of the ones above by using, respectively, Eqs. (C.1), (38), (41), and (51) together with the observed masses of the Higgs boson, top quark, W boson, and the EW vev. The last relation above ensures that the physical W is mostly made of the $SU(2)_L$ gauge field. We scan the parameter space defined above and collect 1000 data points, each satisfying the following constraints,

$$m_{SUSY} > 5 \text{ TeV} , \quad m_{H^0} > 600 \text{ GeV} , \quad m_{\Pi^0}, m_{H^\pm}, m_{h^{\pm\pm}}, m_{H^{\pm\pm}}, m_E, m_N \geq 100 \text{ GeV} , \quad (60)$$

as well as the experimental limits on S and T [39]. We also check that the collected data points satisfy the ATLAS lower limit on the mass of a sequential W' boson [48], equal to 2.55 TeV, after an appropriate rescaling of the limit which takes into account the non-SM value of the W' coupling to fermions [41].

Having at hand a sufficiently large collection of viable points in the 3WT parameter space, we perform a goodness of fit analysis by using the observed Higgs decay rates to $\gamma\gamma$, ZZ , WW , $\tau\tau$, bb , $\gamma\gamma JJ$ at ATLAS [49–53] and CMS [54–57], and to $\gamma\gamma$, WW , and bb at Tevatron [58], as well as the S and T experimental values [39], for a total of 19 observables. The LHC and Tevatron results are expressed in terms of the signal strengths, defined as

$$\hat{\mu}_{ij} = \frac{\sigma_{\text{tot}} \text{Br}_{ij}}{\sigma_{\text{tot}}^{\text{SM}} \text{Br}_{ij}^{\text{SM}}} , \quad \sigma_{\text{tot}} = \sum_{\Phi=h, qqh, \dots} \epsilon_{\Phi} \sigma_{\Phi \rightarrow \Phi'} , \quad \Phi = pp, p\bar{p} , \quad (61)$$

Table 3

Combined signal strengths from LHC and Tevatron experiments.

ij	ATLAS	CMS	Tevatron
ZZ	1.50 ± 0.40	0.91 ± 0.27	
$\gamma\gamma$	1.65 ± 0.32	1.11 ± 0.31	6.20 ± 3.30
WW	1.01 ± 0.31	0.76 ± 0.21	0.89 ± 0.89
$\tau\tau$	0.70 ± 0.70	1.10 ± 0.40	
bb	-0.40 ± 1.10	1.30 ± 0.70	1.54 ± 0.77

Table 4

Signal strengths and efficiencies for Higgs decay to $\gamma\gamma$ associated to a dijet at LHC.

	ATLAS 7 TeV	ATLAS 8 TeV	CMS 7 TeV	CMS 8 TeV
$\gamma\gamma JJ$	2.7 ± 1.9	2.8 ± 1.6	2.9 ± 1.9	0.3 ± 1.3
$pp \rightarrow h$	22.5%	45.0%	26.8%	46.8%
$pp \rightarrow qqh$	76.7%	54.1%	72.5%	51.1%
$pp \rightarrow t\bar{t}h$	0.6%	0.8%	0.6%	1.7%
$pp \rightarrow Vh$	0.1%	0.1%	0%	0.5%

where $\epsilon_{\Phi'}$ is the efficiency associated with the given final state Φ' in an exclusive search, while for inclusive searches one simply has $\sigma_{\text{tot}} = \sigma_{pp \rightarrow h^0(X)}$, the h^0 production total cross section.

The combined signal strengths from ATLAS, CMS,⁷ and Tevatron are given in Table 3, while the signal strengths and efficiencies⁸ for dijet associated $\gamma\gamma$ production at ATLAS and CMS are listed in Table 4.

Finally, the observed values for the S and T parameters are [39]

$$S = 0.04 \pm 0.09, \quad T = 0.07 \pm 0.08, \quad r(S, T) = 88\%, \quad (62)$$

with the last quantity defining the correlation of the two parameters.

For a detailed description of the present fit we refer the reader to [41], where the same statistical analysis has been performed for a different model. Given that no new physics has been detected, only the contributions of new charged particles at one loop to $\Gamma_{h \rightarrow \gamma\gamma}$ become relevant when comparing the 3WT predictions to the data in Tables 3, 4. More explicitly one has [59]

$$\Gamma_{h \rightarrow \gamma\gamma} = \frac{\alpha_e^2 m_h^3}{256\pi^3 v_w^2} \left| \sum_i N_i e_i^2 F_i \right|^2, \quad (63)$$

with i summed over all the charged particles, N_i is the number of colors, e_i the charge in electron units, and F_i a function of the mass m_i and the coupling coefficient defined in [41]. In the limit of new particles being much heavier than the light Higgs, one finds

$$F_{W_i} = 7a_{W_i}, \quad F_E = F_N = -a_f \frac{4}{3}, \quad F_{S_i} = -a_{S_i} \frac{1}{3}, \quad (64)$$

with the coupling coefficients defined by Eq. (52). We can therefore mimic the contribution of the charged non-SM particles in 3WT to the observables in Tables 3, 4 by including only the new contribution of a heavy singly charged vector boson with coupling coefficient $a_{V'}$ determined by

⁷ We use the mass cut based result for CMS result on the Higgs to diphoton decay.

⁸ We chose to include only the loose categories from the ATLAS and CMS dataset at 8 TeV.

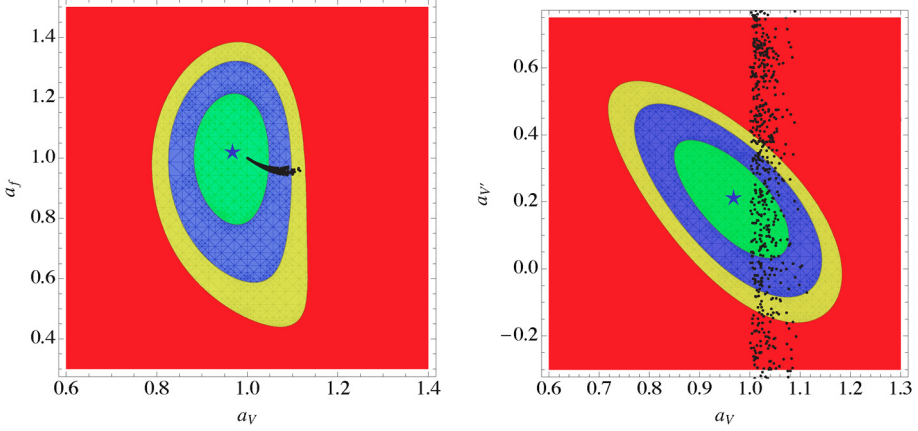


Fig. 1. Viable data points in the (a_V, a_f) (left panel) and $(a_V, a_{V'})$ (right panel) planes, together with the 68% (green), 90% (blue), and 95% (yellow) CL region. The blue star in each plot marks the optimal coupling coefficients on the respective planes. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

$$a_{V'} \equiv \frac{1}{7} (F_{W'} + F_{W''} + 4F_{\Omega}) - \frac{a_S}{21},$$

$$a_S \equiv -3 (16F_E + 4F_N + F_{H\pm} + 4F_{h\pm\pm} + 4F_{H\pm\pm}), \quad (65)$$

where the factors of 4 account for the double charge of the corresponding states. Moreover, to simplify the analysis we redefine consistently with [41]

$$a_Z \approx a_W \equiv a_V, \quad (66)$$

where the numerical deviations from the first approximate equality above turn out to be negligible for the collected data points compared to the uncertainties on the coupling coefficients. At each collected data point we determine the numerical values of a_f , a_V , and $a_{V'}$ by Eqs. (53), (65) and (66), while we calculate numerically the coupling coefficients of the charged scalars. In Fig. 1 we plot the viable data points on the (a_V, a_f) (left panel) and $(a_V, a_{V'})$ (right panel) planes, while in Fig. 2 we plot the data points on the $(a_{V'}, a_f)$ plane, together with the 68% (green), 90% (blue), and 95% (yellow) confidence level (CL) regions. In both plots the missing parameter is fixed to the optimal value marked with a blue star. It is clear from Fig. 1, left panel, that the W and Z couplings are enhanced, compared to their SM values, while the SM fermion couplings are suppressed. This result for the 3WT couplings enhances the Higgs decay to diphotons. However, the contribution of the new charged fermions and scalars, expressed by Eq. (65), is large and interferes destructively with the W contribution to the same process. As a consequence the data point minimizing χ^2 in the (a_f, a_V, a_S) space, obtained from a 3WT particle spectrum without the composite vector resonances at low energy, is ruled out:

$$a_V = 1.00, \quad a_f = 1.00, \quad a_S = 20.5, \quad S = 0.04, \quad T = 0.07;$$

$$\chi_{\min}^2/\text{d.o.f.} = 3.42, \quad P(\chi^2 > \chi_{\min}^2) = 0.0004\%, \quad \text{d.o.f.} = 16. \quad (67)$$

In calculating $\chi_{\min}^2/\text{d.o.f.}$ in the above equations we assumed the model to allow three free parameters (a_f, S, T) , since a_V is strongly correlated with a_f near χ_{\min}^2 and a_S is basically constant. The contribution of the new charged vector bosons, and especially that of the vector baryon Ω , to

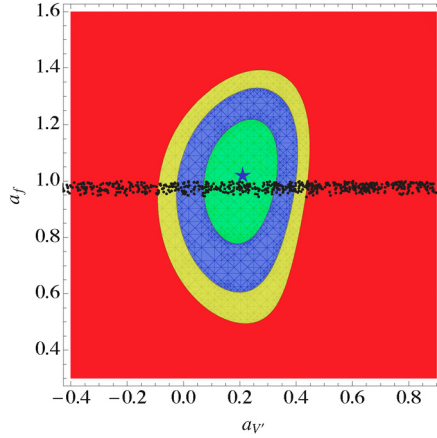


Fig. 2. Viable data points in the $(a_{V'}, a_f)$ plane passing through the point with optimal coupling coefficients in the $(a_{V'}, a_f, a_V)$ space, marked by a blue star, together with the 68% (green), 90% (blue), and 95% (yellow) CL region. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

the Higgs decay into diphoton is large, and offsets entirely the negative contribution of E , N , and charged scalars in Eq. (65). Among the 1000 viable data points, the one producing the minimum value for χ^2 is:

$$a_V = 1.01, \quad a_f = 0.99, \quad a_{V'} = 0.21, \quad S = 0.04, \quad T = 0.07; \\ \chi_{\min}^2/\text{d.o.f.} = 0.83, \quad P(\chi^2 > \chi_{\min}^2) = 65\%, \quad \text{d.o.f.} = 15, \quad (68)$$

where the number of degrees of freedom (d.o.f.) has decreased by one, since $a_{V'}$ is a free parameter. It is interesting to notice that the optimal value of $a_{V'}$ above is equal to the average $a_{V'}$, calculated over the 1000 data points, while the average values of a_f and a_V are, respectively, 0.98 and 1.03, which are very close to the corresponding optimal values given above. This shows that strong dynamics, which we used to determine the scanned range of values of the free parameters, generates rather naturally the coupling strengths favored by LHC data, at least once the direct constraints on the mass spectrum and the EW precision parameters are satisfied.

The 3WT result in Eq. (68) should be compared to the SM one:

$$\chi_{\min}^2/\text{d.o.f.} = 0.89, \quad P(\chi^2 > \chi_{\min}^2) = 60\%, \quad \text{d.o.f.} = 19. \quad (69)$$

While the SM fit is less satisfactory than the 3WT one, it clearly shows that the SM is still perfectly viable in light of present collider data. It is instructive to notice that the fit performed with completely free coupling coefficients, therefore not motivated by any specific underlying theory, produces a worse fit than the 3WT:

$$a_V = 0.97_{-0.11}^{+0.10}, \quad a_f = 1.02_{-0.32}^{+0.25}, \quad a_{V'} = 0.21_{-0.18}^{+0.16}, \\ \chi_{\min}^2/\text{d.o.f.} = 0.85, \quad P(\chi^2 > \chi_{\min}^2) = 62\%, \quad \text{d.o.f.} = 14. \quad (70)$$

This is because the underlying strong dynamics introduces a large correlation between a_f and a_V , hence increasing the number of d.o.f. by one, while achieving a χ_{\min}^2 very close to the corresponding result obtained with free coupling coefficients (Fig. 2).

7. Conclusions

In this paper we derived the low energy effective theory of a supersymmetric model with a new strong interacting sector and tested its viability at the LHC. We started from MSSM extended by a strong interacting $\mathcal{N} = 4$ Super Yang Mills (4SYM) sector as well as by a fourth lepton superfamily. By integrating out all the elementary scalars (as well as gauginos and higgsinos), which we assume to be very heavy, we obtained a Technicolor (TC) sector extended by four-fermion interactions between the (4SYM) TC fermions and the SM ones. Due to these interactions, the TC fermion condensate gives mass to the EW gauge bosons and to the SM fermion as well. The advantage of this setup is twofold: Supersymmetry naturalizes the scalars which allow ETC-type generation of fermion masses, while the strong sector disentangles the SUSY breaking scale from the electroweak scale and solves the little hierarchy problem.

Given that at low energy the strong interacting states form bound states, we constructed the effective Lagrangian at the EW scale expressed in terms of composite scalar and vector fields, in addition to the SM fields. Because the TC potential features an SU(3) global symmetry and the TC coupling is near conformal, we called this model SU(3) walking technicolor (3WT). To test the viability of the model, we worked out the Higgs couplings to the fermion and vector mass eigenstates, as well as the S and T EW parameters. We then scanned the model parameter space for data points featuring a viable mass spectrum, ensuring that the couplings remain perturbative at large scales. By performing a goodness of fit analysis using Higgs physics data from LHC as well as the experimental values of the EW precision parameters, we showed that 3WT fits the experimental data with a level of goodness comparable to that of the SM. Remarkably, the role played by heavy composite vector resonances turned out to be critical, as their contribution to the diphoton decay of the light Higgs is absolutely necessary to bring the corresponding 3WT prediction within the experimental constraints. These composite vector resonances, having mass of $\mathcal{O}(\text{TeV})$, should in principle be observable at LHC.

To conclude, we highlight that SU(3) Walking Technicolor is an UV complete model, which, by avoiding any scalars at the EW scale, in principle solves fine tuning problem. This model, moreover, is favored by Higgs physics and EW precision data at a level comparable to that of the SM.

Acknowledgements

We thank R. Foadi for providing the code to evaluate the EW oblique corrections and for discussions. This work was financially supported by the Academy of Finland project 267842.

Appendix A. EW symmetry breaking in global SU(3) invariant technicolor

The SU(3) symmetry of the microscopic TC Lagrangian is spontaneously broken to the maximal diagonal subgroup, SO(3). The symmetry breaking pattern leaves us with five broken generators with associated Goldstone bosons. Such a breaking is driven by the condensate

$$\langle \eta_i^\alpha \eta_j^\beta \epsilon_{\alpha\beta} E^{ij} \rangle = \langle 2U_R^\dagger U_L + D_L D_L \rangle, \quad (\text{A.1})$$

where the indices $i, j = 1, \dots, 3$ denote the components of the triplet of η , and the Greek indices indicate the ordinary spin. The matrix E is a 3×3 matrix defined as

$$E = \begin{pmatrix} 0 & 0 & 1 \\ 0 & 1 & 0 \\ 1 & 0 & 0 \end{pmatrix}. \tag{A.2}$$

The above condensate is invariant under an SO(3) symmetry. It is convenient to separate the eight generators of SU(3) into the three that leave the vacuum invariant, S^a , and the remaining five that do not, X^a . Then the S^a generators of the SO(3) subgroup satisfy the relation

$$S^a E + E S^{aT} = 0, \quad \text{with} \quad a = 1, \dots, 3, \tag{A.3}$$

so that $uEu^T = E$, for $u \in \text{SO}(3)$. An explicit realization of the generators is shown in [Appendix B](#).

The scalar and pseudoscalar degrees of freedom, necessary to model the Goldstone bosons and spontaneous symmetry breaking, consist of a composite Higgs and its pseudoscalar partner, as well as five pseudoscalar Goldstone bosons and their scalar partners. These can be assembled in the matrix

$$M = \left[\frac{\sigma + i\Theta}{\sqrt{3}} I_3 + \sqrt{2}(i\Pi^a + \tilde{\Pi}^a) X^a \right] E, \tag{A.4}$$

which transforms under the full SU(3) group according to

$$M \rightarrow uMu^T, \quad \text{with} \quad u \in \text{SU}(3). \tag{A.5}$$

The X^a 's, $a = 1, \dots, 5$ are the generators of the SU(3) group which do not leave the vacuum expectation value (VEV) of M invariant

$$\langle M \rangle = \frac{v}{\sqrt{3}} E. \tag{A.6}$$

Appendix B. SU(3) generators

The generators S^i of SO(3) satisfy $S^i E + E S^{iT} = 0$. The other generators of SU(3) are written as X^i . The generators are normalized as

$$\text{Tr}[S^i S^j] = \delta^{ij}/2 \quad \text{Tr}[X^i X^j] = \delta^{ij}/2 \quad \text{Tr}[X^i S^j] = 0 \tag{B.1}$$

and given in terms of the Gell-Mann matrices λ^i by

$$S^1 = \frac{1}{2\sqrt{2}} (\lambda^1 - \lambda^6) \tag{B.2}$$

$$S^2 = \frac{1}{2\sqrt{2}} (\lambda^2 - \lambda^7) \tag{B.3}$$

$$S^3 = \frac{1}{4} (\lambda^3 + \sqrt{3}\lambda^8) \tag{B.4}$$

$$X^1 = \frac{1}{2\sqrt{2}} (\lambda^1 + \lambda^6) \tag{B.5}$$

$$X^2 = \frac{1}{2\sqrt{2}} (\lambda^2 + \lambda^7) \tag{B.6}$$

$$X^3 = \frac{1}{4} (\sqrt{3}\lambda^3 - \lambda^8) \tag{B.7}$$

$$X^4 = \frac{1}{2}\lambda^4 \quad (\text{B.8})$$

$$X^5 = \frac{1}{2}\lambda^5 \quad (\text{B.9})$$

Using the generators above, it is straightforward to obtain the vector and axial-vector charge eigenstates and their elementary particle content from Eqs. (19), (21). First note that the charge operator is $Q = S^3$. We find first the linear combinations of the generators corresponding to charge eigenvalues 0, ± 1 and ± 2 . Then we project the corresponding vector states, e.g. $v_\mu^0 = 2\text{Tr}(S^3 A_\mu)$, and obtain:

$$\begin{aligned} v_\mu^0 &\equiv \frac{A_\mu^3 + \sqrt{3}A_\mu^8}{2} \sim \bar{U}_L \gamma_\mu U_L + \bar{U}_R \gamma_\mu U_R, \\ v_\mu^+ &\equiv \frac{A_\mu^1 - A_\mu^6}{2} - i \frac{A_\mu^2 - A_\mu^7}{2} \sim \bar{D}_L \gamma_\mu U_L + \bar{D}_L^c \gamma_\mu U_R, \\ v_\mu^- &\equiv \frac{A_\mu^1 - A_\mu^6}{2} + i \frac{A_\mu^2 - A_\mu^7}{2} \sim \bar{U}_L \gamma_\mu D_L + \bar{U}_R \gamma_\mu D_L^c, \\ a_\mu^0 &\equiv \frac{\sqrt{3}A_\mu^3 - A_\mu^8}{2} \sim \bar{U}_L \gamma_\mu U_L - \bar{U}_R \gamma_\mu U_R - 2\bar{D}_L \gamma_\mu D_L, \\ a_\mu^+ &\equiv \frac{A_\mu^1 + A_\mu^6}{2} - i \frac{A_\mu^2 + A_\mu^7}{2} \sim \bar{D}_L \gamma_\mu U_L - \bar{D}_L^c \gamma_\mu U_R, \\ a_\mu^- &\equiv \frac{A_\mu^1 + A_\mu^6}{2} + i \frac{A_\mu^2 + A_\mu^7}{2} \sim \bar{U}_L \gamma_\mu D_L - \bar{U}_R \gamma_\mu D_L^c, \\ \Omega_\mu^{++} &\equiv \frac{A_\mu^4 - iA_\mu^5}{\sqrt{2}} \sim \bar{U}_R^c \gamma_\mu U_L, \quad \Omega_\mu^{--} \equiv \frac{A_\mu^4 + iA_\mu^5}{\sqrt{2}} \sim \bar{U}_L \gamma_\mu U_R^c. \end{aligned} \quad (\text{B.10})$$

The particle contents given above reproduce the corresponding results in [33] if one applies there the substitution $D_R \rightarrow D_L^c$.

Appendix C. Squared mass matrices

For the neutral scalar and pseudoscalar states, the charged and doubly charged states, the squared mass matrices are, respectively

$$\mathcal{M}_{\tilde{h}^0}^2 = \begin{pmatrix} 2v_\sigma^2 (\lambda + 2\lambda') & 2v_\Delta v_\sigma (\lambda - 2\lambda'') \\ 2v_\Delta v_\sigma (\lambda - 2\lambda'') & 2(v_\sigma^2 \lambda'' + v_\Delta^2 (\lambda + 4\lambda')) \end{pmatrix}, \quad (\text{C.1})$$

in the $\Re(\sigma^0)$, $\Re(\Delta^0)$ basis,

$$\mathcal{M}_{\tilde{\pi}^0}^2 = \begin{pmatrix} 8v_\Delta^2 \lambda'' & 4v_\Delta v_\sigma \lambda'' \\ 4v_\Delta v_\sigma \lambda'' & 2v_\sigma^2 \lambda'' \end{pmatrix}, \quad (\text{C.2})$$

in the $\Im(\sigma^0)$, $\Im(\Delta^0)$ basis,

$$\mathcal{M}_{\tilde{h}^\pm}^2 = \begin{pmatrix} 2v_\sigma^2 (\lambda'' + \lambda') & 2\sqrt{2}v_\Delta v_\sigma (\lambda'' + \lambda') \\ 2\sqrt{2}v_\Delta v_\sigma (\lambda'' + \lambda') & 4v_\Delta^2 (\lambda'' + \lambda') \end{pmatrix}, \quad (\text{C.3})$$

in the Δ^\pm , σ^\pm basis,

$$\mathcal{M}_{\tilde{h}^{\pm\pm}}^2 = \begin{pmatrix} 2v_\sigma^2\lambda'' - 4(v_\Delta^2 - v_\sigma^2)\lambda' & 4v_\Delta^2\lambda'' + 2v_\sigma^2\lambda' \\ 4v_\Delta^2\lambda'' + 2v_\sigma^2\lambda' & 2v_\sigma^2\lambda'' - 4(v_\Delta^2 - v_\sigma^2)\lambda' \end{pmatrix}, \quad (\text{C.4})$$

in the $\Delta^{\pm\pm}, \delta^{\pm\pm}$ basis.

We define, besides ϵ in Eq. (22), the following dimensionless parameters:

$$x = \frac{g_L v_w}{2m_A}, \quad t_\rho = \frac{\sqrt{2}v_\Delta}{v_\sigma}, \quad z_i = \left(\frac{g_{TC} v_w}{2m_A}\right)^2 r_i, \quad i = 1, 2, 3. \quad (\text{C.5})$$

Then the non-zero terms of the charged vector boson squared mass matrix (which by definition is symmetric) are

$$\begin{aligned} (\mathcal{M}_{\tilde{W}}^2)_{1,1} &= m_A^2 \left[x^2 + \epsilon^2 \left(1 + z_1 - \frac{z_3}{2} (1 + c_\rho^2) \right) \right], \\ (\mathcal{M}_{\tilde{W}}^2)_{2,2} &= m_A^2 \left[1 + z_1 + z_2 s_{2\rho} - \frac{z_3}{2} (1 + c_\rho^2) \right], \quad (\mathcal{M}_{\tilde{W}}^2)_{1,2} = -\frac{\epsilon}{\sqrt{2}} (\mathcal{M}_{\tilde{W}}^2)_{2,2}, \\ (\mathcal{M}_{\tilde{W}}^2)_{3,3} &= m_A^2 \left[1 + z_1 - z_2 s_{2\rho} - \frac{z_3}{2} (1 + c_\rho^2) \right], \quad (\mathcal{M}_{\tilde{W}}^2)_{1,3} = -\frac{\epsilon}{\sqrt{2}} (\mathcal{M}_{\tilde{W}}^2)_{3,3}, \end{aligned} \quad (\text{C.6})$$

in the $\tilde{W}_\mu^\pm, V_\mu^\pm, A_\mu^\pm$ basis, with furthermore the squared mass of the doubly charged vector boson given by

$$m_\Omega^2 = m_A^2 \left[1 + 2c_\rho^2 (z_1 + z_2) - \frac{z_3}{2} (1 + c_\rho^2) \right]. \quad (\text{C.7})$$

Finally, the non-zero terms of the neutral vector boson squared mass matrix in the $\tilde{W}_\mu^3, B_\mu, V_\mu^3, A_\mu^3$ basis are

$$\begin{aligned} (\mathcal{M}_{\tilde{Z}}^2)_{1,1} &= m_A^2 \left[x^2 (1 + s_\rho^2) + \epsilon^2 \left(1 + z_1 - \frac{1}{2} (1 + c_\rho^2) z_3 + z_2 s_\rho^2 \right) \right], \\ (\mathcal{M}_{\tilde{Z}}^2)_{1,2} &= -m_A^2 \left[x^2 (1 + s_\rho^2) + \epsilon^2 (z_1 - c_\rho^2 (2z_1 - z_2) + z_2) \right] t_\xi, \\ (\mathcal{M}_{\tilde{Z}}^2)_{2,2} &= m_A^2 \left[x^2 (1 + s_\rho^2) + \epsilon^2 \left(3 + z_1 + c_\rho^2 (4z_1 - 5z_2) + z_2 - \frac{3}{2} (1 + c_\rho^2) z_3 \right) \right] t_\xi^2, \\ (\mathcal{M}_{\tilde{Z}}^2)_{3,3} &= m_A^2 \left[1 + 2c_\rho^2 (z_1 - z_2) - \frac{z_3}{2} (1 + c_\rho^2) \right], \\ (\mathcal{M}_{\tilde{Z}}^2)_{1,3} &= -\frac{\epsilon}{2} (\mathcal{M}_{\tilde{Z}}^2)_{3,3}, \quad (\mathcal{M}_{\tilde{Z}}^2)_{2,3} = -\frac{3}{2} \epsilon t_\xi (\mathcal{M}_{\tilde{Z}}^2)_{3,3}, \\ (\mathcal{M}_{\tilde{Z}}^2)_{4,4} &= m_A^2 \left[1 + z_1 - z_2 s_{2\rho} - \frac{z_3}{2} (1 + c_\rho^2) \right], \\ (\mathcal{M}_{\tilde{Z}}^2)_{1,4} &= -\frac{\sqrt{3}}{2} \epsilon (\mathcal{M}_{\tilde{Z}}^2)_{4,4}, \quad (\mathcal{M}_{\tilde{Z}}^2)_{2,4} = \frac{\sqrt{3}}{2} \epsilon t_\xi (\mathcal{M}_{\tilde{Z}}^2)_{4,4}. \end{aligned} \quad (\text{C.8})$$

Appendix D. S and T parameters for general neutrino mass matrix

The most general mass terms for a pair of right- and left-handed neutrinos is defined by

$$\mathcal{L} \supset -m_E \bar{E}_R E_L - \frac{1}{2} n_L^T M n_L + \text{h.c.}, \quad M = \begin{pmatrix} M_L & m_D \\ m_D & M_R \end{pmatrix}, \quad n_L = (N_L, \bar{N}_R)^T \quad (\text{D.1})$$

with eigenvalues

$$\lambda_{1,2} = \frac{1}{2} \left[(M_L + M_R) \pm \sqrt{(M_L - M_R)^2 + 4m_D^2} \right]. \quad (\text{D.2})$$

The contributions of the corresponding heavy neutrinos mass eigenstates and of the heavy electron E to the S and T parameters have been derived in terms of integral functions in [46]. From those, we derived the corresponding explicit results:

$$\begin{aligned} S = & \frac{1}{12\pi} \left[1 + 2c_\zeta^4 \left(1 + \log v_1^2 \right) - 2 \log v_E^2 + 2s_\zeta^4 \left(1 + \log v_2^2 \right) \right] \\ & + \frac{s_{2\zeta}^2}{36\pi} \frac{9(1 - \log v_1^2) v_1^4 v_2^2 - 9(1 - \log v_2^2) v_1^2 v_2^4 - (1 - 3 \log v_1^2) v_1^6 + (1 - 3 \log v_2^2) v_2^6}{(v_1^2 - v_2^2)^3} \\ & - \frac{(-1)^\beta s_{2\zeta}^2}{8\pi} \frac{v_1 v_2 \left(v_1^4 - 2v_1^2 v_2^2 \log \frac{v_1^2}{v_2^2} - v_2^4 \right)}{(v_1^2 - v_2^2)^3}, \end{aligned} \quad (\text{D.3})$$

$$\begin{aligned} T = & \frac{\Lambda_{NP}^2}{64\pi c_\zeta^2 s_\zeta^2 m_Z^2} \left[16c_\zeta^4 v_1^2 \log v_1^2 + 16s_\zeta^4 v_2^2 \log v_2^2 + 8v_E^2 \log v_E^2 \right. \\ & - s_{2\zeta}^2 \frac{(1 - 2 \log v_1^2) v_1^4 - (1 - 2 \log v_2^2) v_2^4}{v_1^2 - v_2^2} \\ & + 4(-1)^\beta s_{2\zeta}^2 \frac{(1 - \log v_1^2) v_1^3 v_2 - (1 - \log v_2^2) v_1 v_2^3}{v_1^2 - v_2^2} \\ & + 4c_\zeta^2 \frac{(1 - 2 \log v_1^2) v_1^4 - (1 - 2 \log v_E^2) v_E^4}{v_1^2 - v_E^2} \\ & \left. + 4s_\zeta^2 \frac{(1 - 2 \log v_2^2) v_2^4 - (1 - 2 \log v_E^2) v_E^4}{v_2^2 - v_E^2} \right], \end{aligned} \quad (\text{D.4})$$

where Λ_{NP} is the given renormalization scale, ξ is the EW mixing angle, and

$$\begin{aligned} v_1 = & \frac{\lambda_1}{\Lambda_{NP}}, \quad v_2 = \frac{\lambda_2}{\Lambda_{NP}}, \quad v_E = \frac{m_E}{\Lambda_{NP}}, \quad t_{2\zeta} = \frac{2m_D}{M_R - M_L}, \\ \beta = & \frac{1}{2} \left[1 + \left(\sqrt{\frac{\lambda_1}{|\lambda_1|}} \sqrt{\frac{\lambda_2}{|\lambda_2|}} \right)^2 \right]. \end{aligned} \quad (\text{D.5})$$

In the limit $M_R \rightarrow \infty$, and $M_L = m_E \equiv m_U$, one recovers the results in Eqs. (56).

References

- [1] G. Aad, et al., ATLAS Collaboration, Phys. Lett. B 716 (2012) 1, arXiv:1207.7214.
- [2] S. Chatrchyan, et al., CMS Collaboration, Phys. Lett. B 716 (2012) 30, arXiv:1207.7235.
- [3] The ATLAS Collaboration, Tech. Rep. ATLAS-CONF-2013-034, CERN, Geneva, 2013.
- [4] The CMS Collaboration, Tech. Rep. CMS-PAS-HIG-13-005, CERN, Geneva, 2013.
- [5] M. Heikinheimo, A. Racioppi, M. Raidal, C. Spethmann, K. Tuominen, arXiv:1304.7006, 2013.

- [6] M. Heikinheimo, A. Racioppi, M. Raidal, C. Spethmann, K. Tuominen, arXiv:1305.4182, 2013.
- [7] S. Weinberg, Phys. Rev. D 13 (1976) 974.
- [8] L. Susskind, Phys. Rev. D 20 (1979) 2619.
- [9] D.B. Kaplan, H. Georgi, Phys. Lett. B 136 (1984) 183.
- [10] D.B. Kaplan, H. Georgi, S. Dimopoulos, Phys. Lett. B 136 (1984) 187.
- [11] G. Cacciapaglia, F. Sannino, J. High Energy Phys. 1404 (2014) 111, arXiv:1402.0233.
- [12] K. Yamawaki, M. Bando, K.-i. Matumoto, Phys. Rev. Lett. 56 (1986) 1335.
- [13] M. Bando, K.-i. Matumoto, K. Yamawaki, Phys. Lett. B 178 (1986) 308.
- [14] F. Sannino, K. Tuominen, Phys. Rev. D 71 (2005) 051901, arXiv:hep-ph/0405209.
- [15] D.D. Dietrich, F. Sannino, K. Tuominen, Phys. Rev. D 72 (2005) 055001, arXiv:hep-ph/0505059.
- [16] Z. Chacko, R.K. Mishra, Phys. Rev. D 87 (2013) 115006, arXiv:1209.3022.
- [17] R. Foadi, M.T. Frandsen, F. Sannino, arXiv:1211.1083, 2012.
- [18] S. Di Chiara, R. Foadi, K. Tuominen, Phys. Rev. D 90 (2014) 115016, arXiv:1405.7154.
- [19] S. Di Chiara, R. Foadi, K. Tuominen, S. Tahtinen, arXiv:1412.7835, 2014.
- [20] S. Dimopoulos, L. Susskind, Nucl. Phys. B 155 (1979) 237.
- [21] E. Eichten, K.D. Lane, Phys. Lett. B 90 (1980) 125.
- [22] T. Appelquist, J. Terning, Phys. Rev. D 50 (1994) 2116, arXiv:hep-ph/9311320.
- [23] T. Appelquist, M. Piai, R. Shrock, Phys. Rev. D 69 (2004) 015002, arXiv:hep-ph/0308061.
- [24] E.H. Simmons, Nucl. Phys. B 312 (1989) 253.
- [25] A. Kagan, S. Samuel, Phys. Lett. B 270 (1991) 37.
- [26] C.D. Carone, E.H. Simmons, Nucl. Phys. B 397 (1993) 591, arXiv:hep-ph/9207273.
- [27] C.D. Carone, E.H. Simmons, Y. Su, Phys. Lett. B 344 (1995) 287, arXiv:hep-ph/9410242.
- [28] M. Antola, M. Heikinheimo, F. Sannino, K. Tuominen, J. High Energy Phys. 1003 (2010) 050, arXiv:0910.3681.
- [29] M. Dine, W. Fischler, M. Srednicki, Nucl. Phys. B 189 (1981) 575.
- [30] B.A. Dobrescu, Nucl. Phys. B 449 (1995) 462, arXiv:hep-ph/9504399.
- [31] M. Antola, S. Di Chiara, F. Sannino, K. Tuominen, Eur. Phys. J. C 71 (2011) 1784, arXiv:1001.2040.
- [32] M. Antola, S. Di Chiara, F. Sannino, K. Tuominen, Nucl. Phys. B 864 (2012) 664, arXiv:1111.1009.
- [33] R. Foadi, M.T. Frandsen, T.A. Rytto, F. Sannino, Phys. Rev. D 76 (2007) 055005, arXiv:0706.1696.
- [34] M. Bando, T. Kugo, S. Uehara, K. Yamawaki, T. Yanagida, Phys. Rev. Lett. 54 (1985) 1215.
- [35] M. Bando, T. Kugo, K. Yamawaki, Phys. Rep. 164 (1988) 217.
- [36] H. Georgi, Phys. Lett. B 298 (1993) 187, arXiv:hep-ph/9207278.
- [37] A.G. Cohen, D.B. Kaplan, A.E. Nelson, Phys. Lett. B 412 (1997) 301, arXiv:hep-ph/9706275.
- [38] M. Antola, K. Tuominen, arXiv:1105.3178, 2011.
- [39] J. Beringer, et al., Particle Data Group, Phys. Rev. D 86 (2012) 010001.
- [40] D.D. Dietrich, F. Sannino, Phys. Rev. D 75 (2007) 085018, arXiv:hep-ph/0611341.
- [41] T. Alanne, S. Di Chiara, K. Tuominen, arXiv:1303.3615, 2013.
- [42] M.E. Peskin, T. Takeuchi, Phys. Rev. Lett. 65 (1990) 964.
- [43] M.E. Peskin, T. Takeuchi, Phys. Rev. D 46 (1992) 381.
- [44] R. Barbieri, A. Pomarol, R. Rattazzi, A. Strumia, Nucl. Phys. B 703 (2004) 127, arXiv:hep-ph/0405040.
- [45] R.S. Chivukula, E.H. Simmons, H.-J. He, M. Kurachi, M. Tanabashi, Phys. Rev. D 70 (2004) 075008, arXiv:hep-ph/0406077.
- [46] O. Antipin, M. Heikinheimo, K. Tuominen, J. High Energy Phys. 0910 (2009) 018, arXiv:0905.0622.
- [47] H.-J. He, N. Polonsky, S.-f. Su, Phys. Rev. D 64 (2001) 053004, arXiv:hep-ph/0102144.
- [48] ATLAS-CONF-2012-086, ATLAS-COM-CONF-2012-132 (2012).
- [49] The ATLAS Collaboration, Tech. Rep. ATLAS-CONF-2012-160, CERN, Geneva, 2012.
- [50] The ATLAS Collaboration, Tech. Rep. ATLAS-CONF-2012-161, CERN, Geneva, 2012.
- [51] The ATLAS Collaboration, Tech. Rep. ATLAS-CONF-2013-030, CERN, Geneva, 2013.
- [52] The ATLAS Collaboration, Tech. Rep. ATLAS-CONF-2013-012, CERN, Geneva, 2013.
- [53] The ATLAS Collaboration, Tech. Rep. ATLAS-CONF-2013-013, CERN, Geneva, 2013.
- [54] The CMS Collaboration, Tech. Rep. CMS-PAS-HIG-13-002, CERN, Geneva, 2013.
- [55] The CMS Collaboration, Tech. Rep. CMS-PAS-HIG-13-001, CERN, Geneva, 2013.
- [56] The CMS Collaboration, Tech. Rep. CMS-PAS-HIG-13-003, CERN, Geneva, 2013.
- [57] The CMS Collaboration, Tech. Rep. CMS-PAS-HIG-13-004, CERN, Geneva, 2013.
- [58] T. Aaltonen, et al., CDF Collaboration, D0 Collaboration, arXiv:1303.6346, 2013.
- [59] J.F. Gunion, H.E. Haber, G.L. Kane, S. Dawson, Front. Phys. 80 (2000) 1.