A partially composite Goldstone Higgs

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We consider a model of dynamical electroweak symmetry breaking with a partially composite Goldstone Higgs. The model is based on a strongly-interacting fermionic sector coupled to a fundamental scalar sector via Yukawa interactions. The SU(4) × SU(4) global symmetry of these two sectors is broken to a single SU(4) via Yukawa interactions. Electroweak symmetry breaking is dynamically induced by condensation due to the strong interactions in the new fermionic sector which further breaks the global symmetry SU(4) → Sp(4). The Higgs boson arises as a partially composite state which is an exact Goldstone boson in the limit where SM interactions are turned off. Terms breaking the SU(4) global symmetry explicitly generate a mass for the Goldstone Higgs. The model realizes in different limits both (partially) composite Higgs and (bosonic) Technicolor models, thereby providing a convenient unified framework for phenomenological studies of composite dynamics. It is also a dynamical extension of the recent elementary Goldstone-Higgs model.

Preprint: CP3-Origins-2017-037 DNRF90

I. INTRODUCTION

Gauge-Yukawa models with a strongly interacting fermion sector were proposed in [1] for electroweak symmetry breaking (EWSB) and Standard Model (SM) fermion mass generation. The motivation was to achieve dynamical EWSB and to alleviate the SM naturalness problem, while circumventing the challenges in building a dynamical gauge theory of fermion masses [2–4]. Later developments of this idea include (partially-)Composite-Higgs models (pCH) [5–7] and bosonic Technicolor (bTC) [8–11]. In the former case, the SM-Higgs-like scalar can arise as a mixture between a doublet of Goldstone bosons (GBs) from the composite dynamics and an elementary scalar doublet. In the latter case, a SM-Higgs-like scalar can arise as a mixture of an isosinglet composite resonance and an elementary scalar.

Here we study a model with four Weyl fermions transforming under a new SU(2)TC gauge group and coupled via Yukawa interactions to an SU(4) multiplet of scalars in the two-index antisymmetric representation. This leads to an SU(4)/Sp(4) coset structure and a parameter space encompassing both the (p)CH and (b)TC models while being a dynamical extension of the elementary Goldstone-Higgs model [12].

The SM fermions obtain their masses via ordinary Yukawa couplings to the elementary scalar multiplet. The weak gauge bosons, on the other hand, obtain masses from both a vacuum expectation value (vev) of the elementary Higgs multiplet and the composite sector such that the electroweak (EW) scale $v_w = 246$ GeV is set by

$$v_w^2 = v^2 + f^2 \sin^2 \theta,$$

where $v$ is the vev of the neutral CP-even component of the elementary weak doublet in the SU(4) multiplet, $f$ is the GB decay constant of the composite sector, and $\theta$ is the vacuum-misalignment angle ($\pi/2 \leq \theta \leq \pi$). The TC limit $\theta = \pi/2$, $v = 0$ was studied in e.g. [13, 14] while the CH limit $\pi/2 < \theta < \pi$, $v = 0$ was considered in [15–21]. Finally the pCH limit, $\pi/2 < \theta < \pi$, $v \neq 0$, has recently been studied in [6, 7]. In all cases, the effective low-energy description is based on the SU(4)/Sp(4) coset.

Since the model here is formulated explicitly in terms of elementary constituents, the composite contributions to the spectrum may be predicted using lattice simulations [22, 23]. Differently from [5, 7], the partially composite Higgs remains an exact GB in the presence of the Yukawa interactions between the new fermions and the scalar multiplet. We find a quartic self-coupling of the scalar multiplet which can be larger than in the
SM and help alleviate the potential vacuum stability bounds. The quartic coupling was set to zero in \[6\], and it typically comes out small in bTC models leading to stringent vacuum stability bounds \[11\]. Finally, we also discuss the classically scale-invariant limit of the model.

The paper is organised as follows: In Sec. \[II\] we introduce the SU(4)-symmetric model, and in Sec. \[III\] we discuss the minimal set of explicit breaking sources leading to the desired phenomenology. In Sec. \[IV\] we consider the classically scale-invariant limit, and in Sec. \[V\] we conclude.

II. THE MODEL AND THE EFFECTIVE DESCRIPTION

We consider an underlying SU(2)\(_{\text{TC}}\) TC model with four Weyl fermions in the fundamental representation of SU(2)\(_{\text{TC}}\). The fermion content in terms of left-handed Weyl fields, with \(\bar{U}_L \equiv U_L^\dagger\), along with their EW quantum numbers is presented in Table I.

| SU(2)\(_{\text{TC}}\) | SU(2)\(_{\text{W}}\) | U(1)\(_Y\) |
|-----------------|----------------|----------|
| \((U_L,D_L)\)   | \[\square\]    | \[\square\] 0 |
| \(\bar{U}_L\)   | \[\square\]    | 1        | -1/2 |
| \(\bar{D}_L\)   | \[\square\]    | 1        | +1/2 |

With EW and Yukawa interactions turned off, the new fermion sector features a SU(4) global symmetry, with \(Q_L = (U_L, D_L, \bar{U}_L, \bar{D}_L)\) transforming in the fundamental representation of this SU(4). In addition, we introduce an (elementary) scalar multiplet, \(\Phi\), in the two-index antisymmetric representation of SU(4) transforming as \(\Phi \rightarrow g\Phi g^T\) under SU(4) transformations.

The Lagrangian describing the TC and scalar sectors is given by

\[
\mathcal{L}_{\text{PCH}} = \bar{Q} i \not{D} Q + D_\mu \Phi^\dagger D^\mu \Phi - m_\Phi^2 \Phi^\dagger \Phi - \lambda_\Phi (\Phi^\dagger \Phi)^2 - y_Q Q^T \Phi Q + \text{h.c.,}
\]

where the antisymmetric flavour structure of the fermion bilinear, after antisymmetrising Lorentz and gauge indices, is kept implicit.

This Lagrangian is SU(4) invariant in the limit \(g, g' \rightarrow 0\) and further SU(4)\(_Q\) \(\times\) SU(4)\(_\Phi\) invariant when in addition \(y_Q = 0\) with the two copies of SU(4) acting on the fermionic and scalar sectors, respectively. Furthermore, it is scale invariant at the classical level if \(m_\Phi^2 = 0\).

We first consider the fermionic sector in the limit \(y_Q = 0\). The flavour-antisymmetric condensate \((\epsilon Q_L^T Q_L) \sim f^3 E_I J\) breaks SU(4)\(_Q\) to Sp(4)\(_Q\).

We embed SU(2)\(_L\) \(\times\) SU(2)\(_R\) in SU(4)\(_Q\) by identifying the left and right generators

\[
T_L^i = \frac{1}{2} \left( \sigma_i \ 0 \ 0 \right), \quad T_R^i = \frac{1}{2} \left( 0 \ 0 \ -\sigma_i^T \right), \quad (3)
\]

where \(\sigma_i\) are the Pauli matrices. The EW subgroup is gauged after identifying the generator of hypercharge with \(T_R^3\); see e.g. \[20\] for details.

The breaking of the EW gauge group depends on the alignment between the EW subgroup and the stability group Sp(4)\(_B\). The quartic coupling was set to zero in \[6\], and it typically comes out small in bTC models leading to stringent vacuum stability bounds. The quartic coupling was set to zero in \[6\], and it typically comes out small in bTC models leading to stringent vacuum stability bounds. The quartic coupling was set to zero in \[6\], and it typically comes out small in bTC models leading to stringent vacuum stability bounds.

In the scale-invariant limit of the model, we use the exponential map

\[
\Sigma = \exp \left( \frac{2\sqrt{2} i}{f} \Pi^a \sigma^a \right) E, \quad (5)
\]

where the \(X^a\), with \(a = 1, \ldots, 5\), are the broken generators corresponding to the vacuum \(E\).

The composite Goldstone degrees of freedom in the SU(4)\(_Q\)/Sp(4)\(_B\) coset can be parameterised by the exponential map

\[
\Sigma = \exp \left( \frac{2\sqrt{2} i}{f} \Pi^a \sigma^a \right) E, \quad (5)
\]

where the \(X^a\), with \(a = 1, \ldots, 5\), are the broken generators corresponding to the vacuum \(E\).

We now include the scalar sector and parameterise the scalar multiplet in terms of EW eigenstates according to the vacuum \(E\). To incorporate the EW embedding, we introduce spurions, \(P_a, \tilde{P}_a\),

\[
2P_1 = \delta_1 \delta_{23} - \delta_{13} \delta_{21}, \quad 2P_2 = \delta_2 \delta_{23} - \delta_{13} \delta_{22}, \quad 2\tilde{P}_1 = \delta_{12} \delta_{13} - \delta_{13} \delta_{23}, \quad 2\tilde{P}_2 = \delta_{22} \delta_{13} - \delta_{13} \delta_{23},
\]

such that \(H_a \equiv \text{Tr}[P_a \Phi]\) and \(\tilde{H}_a \equiv \text{Tr}[\tilde{P}_a \Phi]\) transform as EW doublets with hypercharges +1/2 and -1/2, respectively. Furthermore, the projectors to the EW-singlet directions are given by

\[
P^S_1 = \frac{1}{2} \left( -i \sigma_2 \ 0 \ 0 \right), \quad P^S_2 = \frac{1}{2} \left( 0 \ 0 \ i \sigma_2 \right). \quad (7)
\]
In terms of the above projectors, we can write the scalar multiplet as

$$\Phi = \sum_{\alpha=1,2} P_\alpha H_\alpha + \tilde{P}_\alpha \tilde{H}_\alpha + P_1^S S + P_2^S S^*,$$

where

$$H = \frac{1}{\sqrt{2}} \begin{pmatrix} \sigma_h - i\pi_3^h \\ \pi_2^h + i\pi_1^h \end{pmatrix},$$

and

$$S = \frac{1}{\sqrt{2}} (S_R + iS_I)$$

parameterises the EW-singlet scalars.

We couple the SM fermions, in particular the top quark, to the EW-doublet scalar $H$ via the standard Yukawa interactions of the form

$$y q_L \bar{Q} + h.c.$$ where $y \equiv \langle \sigma_h \rangle, v_S \equiv \langle S_R \rangle, \text{ and } m^2_x = m^2_\Phi + \lambda_4 (v^2 + v^2_\delta).$ We use the shorthand notations $s_x \equiv \sin x, c_x \equiv \cos x, t_x \equiv \tan x$ throughout the paper.

The CP-even neutral states $\sigma_h, \Pi_4, S_R$ mix. Using Eq. (1) and

$$t_\beta \equiv \frac{v}{f s_\theta},$$

we can write the mass eigenvalues as

$$m_1^2 = 0, \quad m_2^2 = \frac{m^2_\lambda}{c_\beta^2}, \quad m_3^2 = m^2_\lambda + 2\lambda_4 v^2_\delta s^2_\beta.$$  (17)

The eigenstates corresponding to the eigenvalues are given by

$$h_1 = s_\beta c_\theta \sigma_h + c_\beta \Pi_4 + s_\beta s_\theta S_R, \quad h_2 = c_\beta c_\theta \sigma_h - s_\beta \Pi_4 + c_\beta s_\theta S_R, \quad h_3 = s_\theta \sigma_h - c_\theta S_R.$$  (18)

In the limit of $t_\beta \gg 1$ and $s_\theta \ll 1$ the massless $h_1$ state is mostly $\sigma_h$, the scalar excitation of the elementary doublet. The additional four massless states are the would-be GBs eaten by the $W$ and $Z$ bosons as well as an additional CP-odd state.

III. BREAKING OF THE GLOBAL SU(4) AND A PSEUDO-GOLDSTONE HIGGS

The gauging of the EW subgroup and the Yukawa interactions between $H$ and the SM fermions break the global symmetry explicitly. The dominant EW effect comes from gauge-boson loops with the techniquarks. The one-loop corrections to the elementary scalar potential from the EW and SM-fermion sectors, i.e. the Coleman–Weinberg potential [25], are higher order in perturbation theory. We include the leading EW contribution in the effective potential by adding the effective operators [26, 27]

$$V_{\text{gauge}} = -C_g \left[ g^2 f^4 \sum_{i=1}^3 \text{Tr} (T^i_L \Sigma (T^i_L \Sigma)^*) + g^2 f^4 \text{Tr} (T^3_R \Sigma (T^3_R \Sigma)^*) \right],$$

where $C_g$ is a positive loop-factor, and we expect $C_g \sim O(1)$. For the effective potential, these yield

$$V_{\text{eff}} \supset -\frac{1}{2} \tilde{C}_g z_2^2 f^4 e^2_\theta,$$  (20)

where $\tilde{C}_g \equiv C_g (3g^2 + g^2)$.\]
Another possible source of explicit breaking of the global symmetry is splitting of the masses of the EW-doublet and singlet components of the scalar multiplet, $\Phi$. We expect a small splitting from quantum corrections from the top-quark and EW gauge sectors but here we simply add an explicit mass resulting in such a splitting,

$$V_{\delta m^2} = 2\delta m^2 \text{Tr}[P^S_i \Phi] \text{Tr}[P^S_i \Phi]^* = \frac{1}{2} \delta m^2 (S^2_R + S^2_t).$$

We note that the mass of the fifth CP-odd GB state may be lifted by adding two independent mass terms for the singlets without affecting vacuum alignment.

### A. Vacuum and spectrum

Taking the above sources of explicit SU(4) breaking into account, the vacuum conditions now read (cf. Eq. 15)

$$
y_Q = \frac{m_\chi^2}{8\sqrt{2}\pi Z_2 f^3 s_0},
$$

$$v_S = \frac{\tilde{C}_g Z_2^2 f^2 \beta s_0^2 - v^2 m_\chi^2}{4 g v m_\chi^2},$$

$$\delta m^2 = \frac{\tilde{C}_g Z_2^2 f^2 \beta s_0^2 m_\chi^2}{v^2 m_\chi^2 - \tilde{C}_g Z_2^2 f^2 s_0^2}.$$

The CP-even neutral states $h, \Pi_4, S_R$ again mix, but due to the explicit breaking of SU(4), the lightest mass eigenstate $h_1$ now acquires a nonzero mass and becomes a pseudo-GB (pGB). We fix the mass of this lightest state to the observed Higgs mass, $m_{h_1} = 125$ GeV, and show the values of the Lagrangian mass parameters $f, \delta m^2, m_\phi^2$ for fixed values of $s_0 = 0.1, 0.3$ and $C_g = 1$ in the $(\lambda_\phi, t_\beta)$ plane in Figs. 1 and 2 respectively. Below the cyan line in the upper panel, $m_\phi^2 > 0$, while $\delta m^2 > 0$ throughout the parameter space as shown in the lower panel. In particular, we can achieve correct EWSB driven purely by the strong dynamics $m_\phi^2 > 0, \delta m^2 > 0$. Furthermore, we generically have $\delta m^2 \sim m_\phi^2 \sim f^2$, and for $s_0 \lesssim 0.1$ it is possible to reach scalar masses parameters above a TeV.

The ratio $m_\phi^2/f^2$ can be written for $s_0 \ll 1$ as

$$\frac{m_\phi^2}{f^2} = \frac{8\sqrt{2}\pi Z_2 y_Q}{t_\beta} - \lambda_\phi \left( t_\beta^2 + \frac{\tilde{C}_g Z_2^2}{128\pi^2 y_Q} - \frac{\tilde{C}_g Z_2 t_\beta}{4\sqrt{2}\pi y_Q} \right).$$

If we require $t_\beta > 1.3$ to avoid the Landau pole for the top-Yukawa coupling and the flavour-changing neutral currents from the heavy pseudoscalars and restrict $\lambda_\phi < 1$, then the Higgs mass constraint $m_{h_1} = 125$ GeV results in $y_Q \lesssim 0.1$ and further $m_\phi^2/f^2 \lesssim 4$ with the maximum at large $\lambda_\phi$ and small $t_\beta$.

The masses of the heavy pion triplet can be written as

$$m^2_{\pi} = \frac{8\sqrt{2}\pi Z_2 y_Q^2}{t_\beta \gamma^2},$$

and they are of the same order as the heavy scalar masses. The additional scalars are also relatively heavy scaling roughly with $\delta m^2$. The two remaining CP-odd states in the spectrum are the mass eigenstates composed of $\Pi_5$ and $S_I$.

### B. Couplings

To study the viability of the model in light of the current experimental data, we parameterize the rotation to the mass eigenbasis by

$$\left( \begin{array}{c} h_1 \\ h_2 \\ h_3 \end{array} \right) = R \left( \begin{array}{c} \sigma_h \\ \Pi_4 \\ S_R \end{array} \right),$$

and define the coefficients

$$\kappa_{t} \equiv \frac{g_{hIt}^{SM}}{g_{hIt}^{\tilde{g}}} = \frac{y_{t} R_{11}}{y_{t}^{SM}} = \frac{R_{11}}{s_{\beta}},$$

and

$$\kappa_{V} \equiv \frac{g_{hW+W-}^{SM}}{g_{hW+W-}^{\tilde{g}}} = R_{11} s_{\beta} + R_{12} c_{\beta} s_{\beta}.$$

We solve the $R$ matrix numerically and show the values of $\kappa_t$-coefficient for fixed values of $s_0 = 0.1$ and $C_g = 1$ in the $(\lambda_\phi, t_\beta)$ plane in Fig. 3; the values of $\kappa_V$ are almost identical, and therefore we do not show them separately here.

The LHC experiments constrain the couplings of the lightest scalar eigenstate, the Higgs boson, to top quark and EW gauge bosons. The combined ATLAS & CMS analysis of the Run-1 data [25] limits the modifications of the vector and fermion couplings in a two-parameter fit of $\kappa_f = \kappa_t$ and $\kappa_V$ to less than 20% at 1σ level. It is evident that in this respect the model is viable in most of the parameter space.

### IV. THE CLASSICALLY SCALE-INARIANT LIMIT

Figs. 1 and 2 show that it is possible to have $m_\phi^2 = 0$, as along the cyan line. In the limit
where both $\delta m^2 = 0$, the Lagrangian in Eq. (2) then becomes scale invariant at the classical level. Even though the scale invariance is broken at the quantum level, the classical scale invariance has been invoked as a guiding principle for the EWSB sector. The model here is distinct from models [25, 29–32] relying on the Coleman–Weinberg mechanism [25] to generate EWSB via one-loop corrections to the elementary scalar potential from the EW and SM-fermion sectors.

Here, EWSB is again induced and communicated to the elementary scalars due to condensation of the technifermions. This is similar to the scale-invariant model of [33, 34] where a singlet scalar is coupled both via Yukawa interactions to new strongly-interacting fermions without SM charges and via quartic couplings to the Higgs.

The corrections from the EW gauge-boson loops to the techniquarks are again important, and we include these contributions to the effective potential as in Eq. (19).

To obtain a non-zero mass for the pGB Higgs, we need to add a source of explicit global symmetry breaking. Contrary to the previous discussion, in the scale-invariant framework we add a splitting between the singlet and doublet quartic couplings instead of mass splitting. This is achieved by adding to the effective potential a term

$$V_{\delta \lambda} = \delta \lambda \text{Tr} \left[ \Phi_S^\dagger \Phi_S \right]^2,$$

(28)
FIG. 3. The (absolute) value of the coefficient $|\kappa_t|$ is depicted for a fixed values of $s_\theta = 0.1$ and $C_g = 1$. The mass of the lightest scalar eigenstate is fixed to 125 GeV. On the gray shaded region on the left, no solution for the Higgs mass condition is found.

where $\Phi_S = P^S S + P^S S^*$. Due to the different RG running of the EW-singlet and -doublet parts, some splitting of the quartics is anticipated, and we also note that the RG running would further induce splitting of the form $\text{Tr} [H^\dagger H] \text{Tr} \left[ \Phi_S^\dagger \Phi_S \right]$. Here we remain agnostic about the origin of the splitting and consider the minimal scenario in Eq. (28).

Minimising the potential then leads to

$$y_Q = \frac{\tilde{C}_g Z_2 f s_2 \theta}{16\sqrt{2}\pi (v c_\theta + v s_\theta)},$$

$$\lambda_\Phi = \frac{\tilde{C}_g Z_2^2 f^4 s_2 \theta}{2v (v^2 + v_S^2)} (v c_\theta + v_S s_\theta),$$

$$\delta \lambda = - \frac{\tilde{C}_g Z_2^2 f^4 s_2 \theta}{2v v_S^3}.$$

Fixing again the mass of the lightest neutral scalar eigenstate to 125 GeV, we show the values of the scalar quartics, and the Yukawa coupling fixing $C_g = 1$ in the $(s_\theta, t_\beta)$ plane in Fig. 4. We disregard the grey areas where the quartic couplings are larger than one, and we again find viable parameter space for a range of $t_\beta$ and $s_\theta$ values. For relatively small $t_\beta$, it is possible to have $\delta \lambda < \lambda_\Phi < 1$. As in the previous case, the Yukawa couplings must be small at the $O(10^{-2} \ldots 10^{-1})$ level.

We leave a more detailed study of the phenomenology of the model presented here to future work. We also note that additional source of SU(4) breaking is needed to lift the mass of the remaining CP-odd GB. Minimally this can be achieved by adding a term of the form $\delta \lambda_S S_R^i S_I^j$. Adding such a term does not affect the previous results, and we leave a detailed investigation of the heavy scalar spectrum for future studies.

V. CONCLUSIONS

In this paper we have proposed a model of a partially composite Goldstone Higgs. The model combines dynamical EWSB with SM fermion mass generation via ordinary Yukawa couplings to a scalar multiplet. The model features a spontaneously broken SU(4) global symmetry with the Higgs as a Goldstone boson in the limit of no SM interactions. The Higgs state is mostly elemen-
tary, and thus its couplings to both the SM-vector and -fermion states are SM Higgs like, while we find a Higgs self-coupling that can be larger than in the SM and thereby help alleviate the vacuum stability constraints. This is in contrast to the fully elementary realisation, where the quartic coupling is expected to be smaller than in the SM [12], thereby providing a potential diagnostics to distinguish these different realisations in future collider experiments.

The SU(4) symmetry is broken by the EW interactions and by adding explicit breaking terms to the scalar potential. We find a viable solution with a SM-like Higgs state from an SU(4)-breaking mass term which is small compared to the compositeness scale $4\pi f$. The SU(4)-preserving mass parameter $m_\Phi^2$ remains of the order of $f^2$, but can be positive as opposed to in the SM, and the symmetry breaking can be fully dynamical. We also find a viable solution in the classically scale-invariant limit by setting $m_\Phi^2 = 0$ and adding SU(4)-non-invariant quartic couplings to the scalar potential. In this case the SU(4)-breaking quartic coupling may be smaller than the SU(4)-preserving one.

In all cases, the additional SU(4) breaking (on top of the SM gauge and Yukawa interactions) can be attained on the scalar-potential level, while the new fermion sector remains SU(4) symmetric.

**ACKNOWLEDGMENTS**

We thank M. Jørgensen, M.L.A. Kristensen, C. Pica and K. Tuominen for discussions. TA acknowledges partial funding from a Villum foundation grant when part of this article was being completed. MTF acknowledges partial funding from The Council For Independent Research, grant number DFF 6108-00623. The CP3-Origins centre is partially funded by the Danish National Research Foundation, grant number DNRF90.