SDU&

University of Southern Denmark

Partially composite Goldstone Higgs boson

Alanne, Tommi; Franzosi, Diogo Buarque; Frandsen, Mads T.

Go to publication entry in University of Southern Denmark's Research Portal

Partially composite Goldstone Higgs boson

Tommi Alanne[*]

Max-Planck-Institut für Kernphysik, Saupfercheckweg 1, 69117 Heidelberg, Germany

Diogo Buarque Franzosi[†]

II. Physikalisches Institut, Universität Göttingen, Friedrich-Hund-Platz 1, 37077 Göttingen, Germany

Mads T. Frandsen[3‡]

CP³-Origins, University of Southern Denmark, Campusvej 55, DK-5230 Odense M, Denmark

(Received 5 October 2017; published 16 November 2017)

We consider a model of dynamical electroweak symmetry breaking with a partially composite Goldstone Higgs boson. The model is based on a strongly interacting fermionic sector coupled to a fundamental scalar sector via Yukawa interactions. The $SU(4) \times SU(4)$ 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 $SU(4) \to Sp(4)$. 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 boson. 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.

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) [5–10] and bosonic technicolor (bTC) [11–14]. 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. Condensation of the technifermions leads to dynamical EWSB and induces a vacuum expectation value (vev) for the scalar multiplet. 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 [15–18].

The SM fermions obtain their Dirac masses via ordinary Yukawa couplings to the elementary scalar multiplet, while the technifermions also obtain vectorlike masses. The weak gauge bosons, on the other hand, obtain masses from both a vev of the elementary Higgs multiplet and the composite sector such that the electroweak (EW) scale $v_{\mathrm{w}} = 246$ GeV is set by

$$v_{\mathrm{w}}^2 = v^2 + f^2 \sin^2 \theta, \qquad (1)$$

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 θ is the vacuum-misalignment angle ($\pi/2 \leq \theta \leq \pi$). The TC limit $\theta = \pi/2$, $v = 0$ was studied in e.g. [19,20], while the CH limit $\pi/2 < \theta < \pi$, $v = 0$ was considered in [21–27]. Finally the pCH limit, $\pi/2 < \theta < \pi$, $v \neq 0$, has recently been studied in [9,10]. 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 [28,29]. Differently from [5,9,10], 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 that in the SM and helps

[*]tommi.alanne@mpi-hd.mpg.de
[†]dbuarqu@gwdg.de
[‡]frandsen@cp3.sdu.dk

alleviate the potential vacuum stability bounds. The quartic coupling was set to zero in [9], and it typically comes out small in bTC models, leading to stringent vacuum stability bounds [14] unless the model is further extended, e.g. into a supersymmetric framework. Finally, we also discuss the classically scale-invariant limit of the model.

The paper is organized 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 $SU(2)_{TC}$ TC model with four Weyl fermions in the fundamental representation of $SU(2)_{TC}$. The fermion content in terms of left-handed Weyl fields, with $\tilde{U}_L \equiv \epsilon U_R^*$, along with their EW quantum numbers is presented in Table I.

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

The Lagrangian describing the TC and scalar sectors is given by

$$\mathcal{L}_{PCH} = \bar{Q}i\slashed{D}Q$$
$$+ D_\mu\Phi^\dagger D^\mu\Phi - m_\Phi^2\Phi^\dagger\Phi - \lambda_\Phi(\Phi^\dagger\Phi)^2$$
$$- y_Q Q^T\Phi Q + \text{H.c.}, \qquad (2)$$

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

This Lagrangian is SU(4) invariant in the limit $g, g' \to 0$ 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 flavor-antisymmetric condensate $\langle \epsilon Q_I^T Q_J \rangle \sim f^3 E_{IJ}$ 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,

TABLE I. The new fermion content.

	$SU(2)_{TC}$	$SU(2)_W$	$U(1)_Y$
(U_L, D_L)	□	□	0
$\tilde{U}_L$	□	1	$-1/2$
$\tilde{D}_L$	□	1	$+1/2$

$$T_L^i = \frac{1}{2}\begin{pmatrix} \sigma_i & 0 \\ 0 & 0 \end{pmatrix}, \qquad T_R^i = \frac{1}{2}\begin{pmatrix} 0 & 0 \\ 0 & -\sigma_i^T \end{pmatrix}, \quad (3)$$

where σ_i are the Pauli matrices. The EW subgroup is gauged after identifying the generator of hypercharge with T_R^3; see e.g. [26,27,30] for details.

The breaking of the EW gauge group depends on the alignment between the EW subgroup and the stability group $Sp(4)_Q$. This can be conveniently parametrized by an angle, θ, after identifying the vacua that leave the EW symmetry intact, $E_\pm$, and the one, E_B, breaking it completely to $U(1)_{EM}$ of electromagnetism. The true vacuum can be written as a linear combination of the EW preserving and breaking vacua, $E = \cos\theta E_- + \sin\theta E_B$. Either choice of $E_\pm$ is equivalent, and here we have chosen E_-. These different vacua can be written explicitly [23] as

$$E_\pm = \begin{pmatrix} i\sigma_2 & 0 \\ 0 & \pm i\sigma_2 \end{pmatrix}, \qquad E_B = \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}, \quad (4)$$

and the angle θ is that in Eq. (1).

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

$$\Sigma = \exp\left(\frac{2\sqrt{2}i}{f}\Pi^a X^a\right)E, \qquad (5)$$

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 parametrize the scalar multiplet in terms of EW eigenstates according to the vacuum E_-. To incorporate the EW embedding, we introduce spurions, P_α, $\tilde{P}_\alpha$,

$$2P_1 = \delta_{i1}\delta_{j3} - \delta_{i3}\delta_{j1}, \qquad 2P_2 = \delta_{i2}\delta_{j3} - \delta_{i3}\delta_{j2},$$
$$2\tilde{P}_1 = \delta_{i1}\delta_{j4} - \delta_{i4}\delta_{j1}, \qquad 2\tilde{P}_2 = \delta_{i2}\delta_{j4} - \delta_{i4}\delta_{j2}, \qquad (6)$$

such that $H_\alpha \equiv \text{Tr}[P_\alpha\Phi]$ and $\tilde{H}_\alpha \equiv \text{Tr}[\tilde{P}_\alpha\Phi]$ transform as EW doublets with hypercharges $+1/2$ and $-1/2$, respectively. Furthermore, the projectors to the EW-singlet directions are given by

$$P_1^S = \frac{1}{2}\begin{pmatrix} -i\sigma_2 & 0 \\ 0 & 0 \end{pmatrix}, \qquad P_2^S = \frac{1}{2}\begin{pmatrix} 0 & 0 \\ 0 & i\sigma_2 \end{pmatrix}. \quad (7)$$

In terms of the above projectors, we can write the scalar multiplet as

$$\Phi = \sum_{\alpha=1,2} P_\alpha H_\alpha + \tilde{P}_\alpha\tilde{H}_\alpha + P_1^S S + P_2^S S^*, \qquad (8)$$

where

$$H = \frac{1}{\sqrt{2}} \begin{pmatrix} \sigma_h - i\pi_h^3 \\ -(\pi_h^2 + i\pi_h^1) \end{pmatrix}, \qquad (9)$$

and $S = \frac{1}{\sqrt{2}}(S_R + iS_I)$ parametrizes 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 $y_t \bar{q}_L H t_R + \text{H.c.}$ where $q_L = (t_L, b_L)$. The techniquarks also couple to the singlet scalar, S, via the Yukawa interacton in Eq. (2) which induces a vectorlike mass term for them.

Turning on y_Q breaks the $SU(4)_Q \times SU(4)_\Phi \to SU(4)$, and upon condensation of the technifermions the $SU(4)$ spontaneously breaks to $Sp(4)$. The five GBs may be classified as an EW-doublet Goldstone Higgs and an EW singlet.

Below the condensation scale, the Lagrangian of Eq. (2) yields

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{kin}} - V_{\text{eff}}, \qquad (10)$$

where, after gauging the EW subgroup, the kinetic terms read

$$\mathcal{L}_{\text{kin}} = \text{Tr}[D_\mu \Phi^\dagger D^\mu \Phi] + \frac{f^2}{8} \text{Tr}[D_\mu \Sigma^\dagger D^\mu \Sigma], \qquad (11)$$

with

$$D_\mu X = \partial_\mu X - i(G_\mu X + X G_\mu^T), \qquad (12)$$

where $X = \Phi, \Sigma$, and the EW gauge fields are encoded in

$$G_\mu = g W_\mu^i T_L^i + g' B_\mu T_R^3. \qquad (13)$$

Furthermore, the effective tree-level potential is given by

$$\begin{aligned} V_{\text{eff}} &= 4\pi f^3 Z_2(y_Q \text{Tr}[\Phi\Sigma] + \text{H.c.}) \\ &\quad + m_\Phi^2 \text{Tr}[\Phi^\dagger \Phi] + \lambda_\Phi \text{Tr}[\Phi^\dagger \Phi]^2, \end{aligned} \qquad (14)$$

where Z_2 is a nonperturbative constant expected to be $O(1)$. We use the numerical value $Z_2 \approx 1.5$ following from the determination of the condensate, in the $SU(2)$ gauge theory with two flavors, in [28].

Then the vacuum conditions read

$$y_Q = \frac{m_\lambda^2 v}{8\sqrt{2}\pi Z_2 f^3 s_\theta}, \qquad t_\theta = -\frac{v}{v_S}, \qquad (15)$$

where $v \equiv \langle \sigma_h \rangle$, $v_S \equiv \langle S_R \rangle$, and $m_\lambda^2 \equiv m_\Phi^2 + \lambda_\Phi(v^2 + v_S^2)$. We use the shorthand notations $s_x \equiv \sin x$, $c_x \equiv \cos x$, and $t_x \equiv \tan x$ throughout the paper.

The CP-even neutral states σ_h, Π_4, and S_R mix. Using Eq. (1) and

$$t_\beta \equiv \frac{v}{f s_\theta}, \qquad (16)$$

we can write the mass eigenvalues as

$$m_1^2 = 0, \qquad m_2^2 = \frac{m_\lambda^2}{c_\beta^2}, \qquad m_3^2 = m_\lambda^2 + 2\lambda_\Phi v_w^2 \frac{s_\beta^2}{s_\theta^2}. \qquad (17)$$

The eigenstates corresponding to the eigenvalues are given by

$$\begin{aligned} h_1 &= s_\beta c_\theta \sigma_h + c_\beta \Pi_4 + s_\beta s_\theta S_R, \\ h_2 &= c_\beta c_\theta \sigma_h - s_\beta \Pi_4 + c_\beta s_\theta S_R, \\ h_3 &= s_\theta \sigma_h - c_\theta S_R. \end{aligned} \qquad (18)$$

In the limit of $t_\beta \gg 1$ and $s_\theta \ll 1$ the massless h_1 state is mostly σ_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 [31], are higher order in perturbation theory. We include the leading EW contribution in the effective potential by adding the effective operators [32,33],

$$\begin{aligned} V_{\text{gauge}} &= -C_g \Big[g^2 f^4 \sum_{i=1}^3 \text{Tr}(T_L^i \Sigma(T_L^i \Sigma)^*) \\ &\quad + g'^2 f^4 \text{Tr}(T_R^3 \Sigma(T_R^3 \Sigma)^*) \Big], \end{aligned} \qquad (19)$$

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

$$V_{\text{eff}} \supset -\frac{1}{2}\tilde{C}_g Z_2^2 f^4 c_\theta^2, \qquad (20)$$

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, Φ. 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,

$$\begin{aligned} V_{\delta m^2} &= 2\delta m^2 \text{Tr}[P_i^S \Phi] \text{Tr}[P_i^S \Phi]^* \\ &= \frac{1}{2}\delta m^2(S_R^2 + S_I^2). \end{aligned} \qquad (21)$$

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)]

$$y_Q = \frac{m_\lambda^2 v}{8\sqrt{2}\pi Z_2 f^3 s_\theta},$$

$$v_S = \frac{\tilde{C}_g Z_2^2 f^4 s_\theta^2 - v^2 m_\lambda^2}{t_\theta v m_\lambda^2},$$

$$\delta m^2 = \frac{\tilde{C}_g Z_2^2 f^4 s_\theta^2 m_\lambda^2}{v^2 m_\lambda^2 - \tilde{C}_g Z_2^2 f^4 s_\theta^2}. \tag{22}$$

The *CP*-even neutral states σ_h, Π_4, and 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, δm^2, and m_Φ^2 for fixed values of $s_\theta = 0.1$, 0.3 and $C_g = 1$ in the (λ_Φ, t_β) 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 $s_\theta \lesssim 0.1$ it is possible to reach scalar mass parameters above a TeV.

The ratio m_Φ^2/f^2 can be written for $s_\theta \ll 1$ as

$$\frac{m_\Phi^2}{f^2} = \frac{8\sqrt{2}\pi Z_2 y_Q}{t_\beta} - \lambda_\Phi \left(t_\beta^2 + \frac{\tilde{C}_g^2 Z_2^2}{128\pi^2 y_Q^2} - \frac{\tilde{C}_g Z_2 t_\beta}{4\sqrt{2}\pi y_Q} \right). \tag{23}$$

If we require $t_\beta > 1.3$ to avoid the Landau pole for the top-Yukawa coupling and the flavor-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 λ_Φ and small t_β.

The masses of the heavy pion triplet can be written as

$$m_\pi^2 = \frac{8\sqrt{2}\pi Z_2 v_w^2 y_Q}{t_\beta s_\theta^2}, \tag{24}$$

and they are of the same order as the heavy scalar masses. The additional scalars are also relatively heavy, scaling roughly with δm^2. The two remaining *CP*-odd states in the spectrum are the mass eigenstates composed of Π_5 and S_I.

B. Couplings

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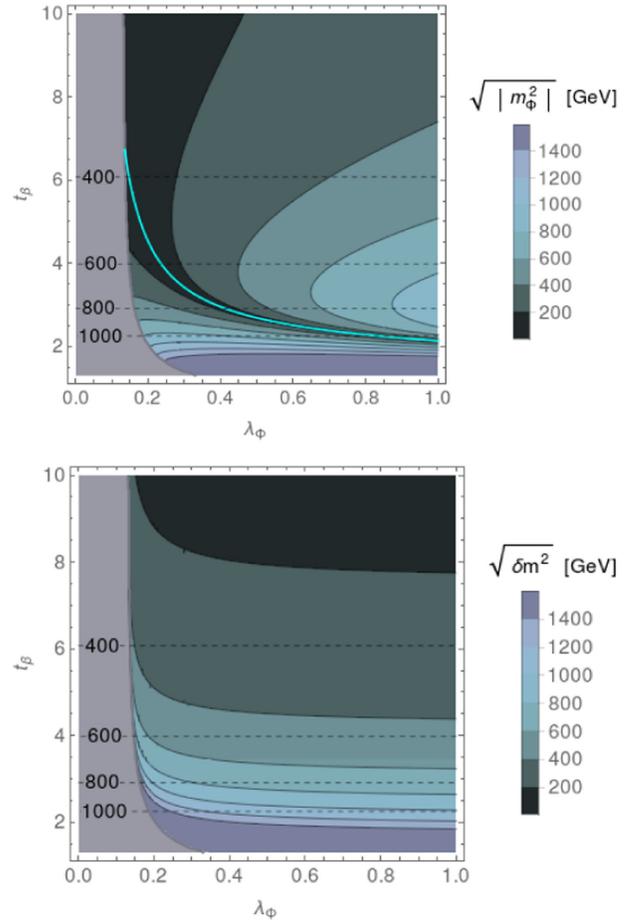

FIG. 1. Contours of the absolute value of the mass parameter $\sqrt{|m_\Phi^2|}$ ($\sqrt{\delta m^2}$) are depicted in the upper (lower) panel for fixed values of $s_\theta = 0.1$ and $C_g = 1$. The mass of the lightest scalar eigenstate is fixed to 125 GeV. On the gray shaded region on the left, no solution for the Higgs mass condition is found. In the upper panel, below the solid cyan curve, $m_\Phi^2 > 0$. In the lower panel, δm^2 is always positive. Dashed lines indicate the value of the composite-pion decay constant, f, in GeV.

$$\begin{pmatrix} h_1 \\ h_2 \\ h_3 \end{pmatrix} = R \begin{pmatrix} \sigma_h \\ \Pi_4 \\ S_R \end{pmatrix}, \tag{25}$$

and define the coefficients

$$\kappa_t \equiv \frac{g_{h_1 \bar{t} t}}{g_{h \bar{t} t}^{\text{SM}}} = \frac{y_t R_{11}}{y_t^{\text{SM}}} = \frac{R_{11}}{s_\beta}, \tag{26}$$

and

$$\kappa_V \equiv \frac{g_{h_1 W^+ W^-}}{g_{h W^+ W^-}^{\text{SM}}} = R_{11} s_\beta + R_{12} c_\beta c_\theta. \tag{27}$$

We solve the R matrix numerically and show the values of the κ_t coefficient for fixed values of $s_\theta = 0.1$ and $C_g = 1$ in the (λ_Φ, t_β) plane in Fig. 3; the values of κ_V are almost identical, and therefore we do not show them separately here.
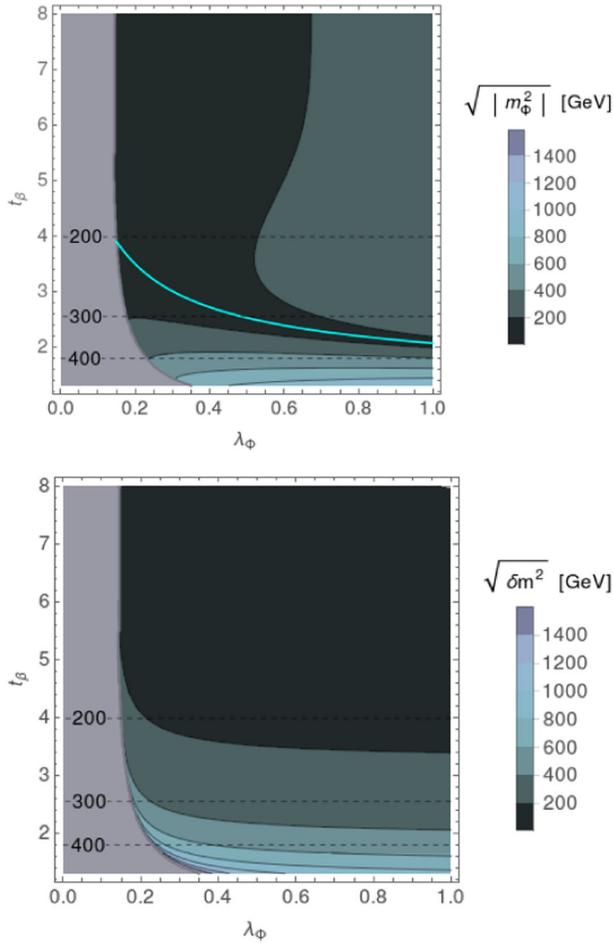
FIG. 2. Contours of the absolute value of the mass parameter $\sqrt{|m_\Phi^2|}$ ($\sqrt{\delta m^2}$) are depicted in the upper (lower) panel for fixed values of $s_\theta = 0.3$ and $C_g = 1$. The mass of the lightest scalar eigenstate h_1 is fixed to 125 GeV. On the gray shaded region on the left, no solution for the Higgs mass condition is found. In the upper panel, below the solid cyan curve, $m_\Phi^2 > 0$. In the lower panel, δm^2 is always positive. Dashed lines indicate the value of the composite-pion decay constant, f, in GeV.

The LHC experiments constrain the couplings of the lightest scalar eigenstate, the Higgs boson, to top quark and EW gauge bosons. The combined ATLAS and CMS analysis of the Run-1 data [34] limits the modifications of the vector and fermion couplings in a two-parameter fit of $\kappa_f = \kappa_t$ and κ_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-INVARIANT LIMIT

Figures 1 and 2 show that it is possible to have $m_\Phi^2 = 0$, as along the cyan line. In the limit where both $m_\Phi^2 = \delta m^2 = 0$, the Lagrangian in Eq. (2) then becomes scale invariant at the classical level. Even though the scale invariance is broken at
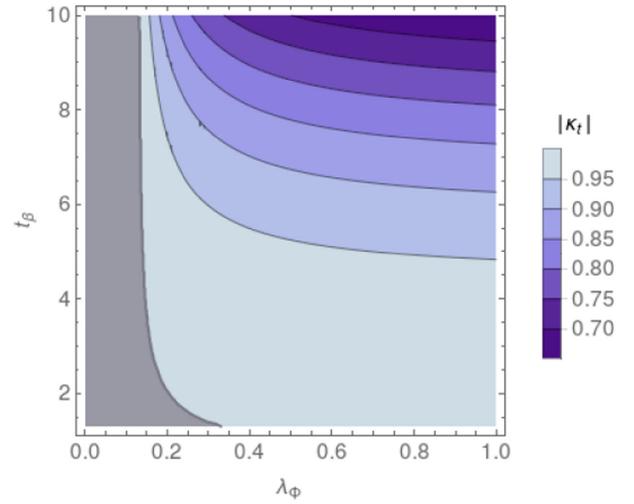

FIG. 3. The (absolute) value of the coeffcient $|\kappa_t|$ is depicted for fixed values of $s_\theta = 0.1$ and $C_g = 1$. The mass of the lightest scalar eigenstate is fixed to 125 GeV. On the gray shaded region on the left, no solution for the Higgs mass condition is found.

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 [31,35–38] relying on the Coleman-Weinberg mechanism [31] 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 [39,40] 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 nonzero 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

$$V_{\delta\lambda} = \delta\lambda \text{Tr}[\Phi_S^\dagger \Phi_S]^2, \qquad (28)$$

where $\Phi_S = P_1^S S + P_2^S S^*$. 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 $\text{Tr}[H^\dagger H]\text{Tr}[\Phi_S^\dagger \Phi_S]$. Here we remain agnostic about the origin of the splitting and consider the minimal scenario in Eq. (28).
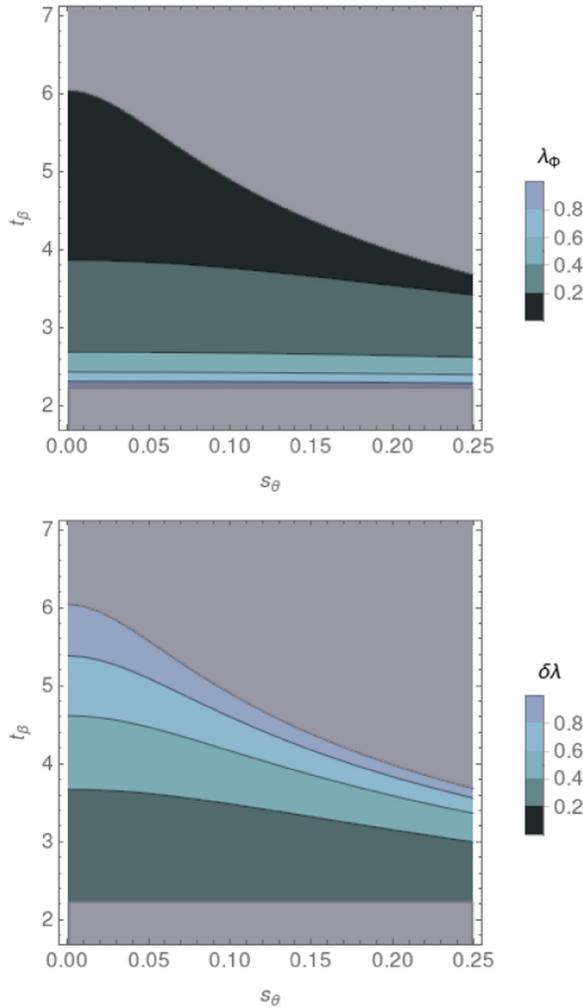
We leave a more detailed study of the phenomenology of the model presented here to future work. We also note that an additional source of SU(4) breaking is needed to lift the mass of the remaining *CP*-odd GB. Minimally this can be achieved by adding a term of the form $\delta\lambda_S S_R^2 S_I^2$. 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 boson. 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 elementary, 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 that in the SM and thereby help alleviate the vacuum stability constraints. This is in contrast to the fully elementary realization, where the quartic coupling is expected to be smaller than that in the SM [16], thereby providing potential diagnostics to distinguish these different realizations 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_Φ^2 remains on the order of f^2 but can be positive as opposed to that 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)-noninvariant 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.

FIG. 4. Values of the SU(4)-invariant quartic coupling, λ_Φ (upper panel), and the SU(4)-breaking one, $\delta\lambda$ (lower panel), for a fixed value $C_g = 1$. The mass of the lightest scalar eigenstate, h_1, is fixed to 125 GeV. On the gray shaded regions, one of the couplings becomes large ($\gtrsim 1$).

Minimizing the potential then leads to

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

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

$$\delta\lambda = -\frac{\tilde{C}_g Z_2^2 f^4 s_{2\theta}}{2vv_S^3}. \tag{29}$$

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_θ, t_β) plane in Fig. 4. We disregard the gray areas where the quartic couplings are larger than one, and we again find viable parameter space for a range of t_β and s_θ values. For relatively small t_β, it is possible to have $\delta\lambda < \lambda_\Phi < 1$. As in the previous case, the Yukawa couplings must be small at the $O(10^{-2}\ldots10^{-1})$ level.

[1] G. 't Hooft, NATO Sci. Ser. B **59**, 135 (1980).
[2] E. Eichten and K. D. Lane, Phys. Lett. **90B**, 125 (1980).
[3] S. Dimopoulos and L. Susskind, Nucl. Phys. **B155**, 237 (1979).
[4] D. B. Kaplan, Nucl. Phys. **B365**, 259 (1991).
[5] D. B. Kaplan and H. Georgi, Phys. Lett. **136B**, 183 (1984).
[6] D. B. Kaplan, H. Georgi, and S. Dimopoulos, Phys. Lett. **136B**, 187 (1984).
[7] T. Banks, Nucl. Phys. **B243**, 125 (1984).
[8] M. J. Dugan, H. Georgi, and D. B. Kaplan, Nucl. Phys. **B254**, 299 (1985).
[9] J. Galloway, A. L. Kagan, and A. Martin, Phys. Rev. D **95**, 035038 (2017).
[10] A. Agugliaro, O. Antipin, D. Becciolini, S. De Curtis, and M. Redi, Phys. Rev. D **95**, 035019 (2017).
[11] E. H. Simmons, Nucl. Phys. **B312**, 253 (1989).
[12] S. Samuel, Nucl. Phys. **B347**, 625 (1990).
[13] A. Kagan and S. Samuel, Phys. Lett. B **252**, 605 (1990).
[14] C. D. Carone, Phys. Rev. D **86**, 055011 (2012).
[15] R. Barbieri, B. Bellazzini, V. S. Rychkov, and A. Varagnolo, Phys. Rev. D **76**, 115008 (2007).
[16] T. Alanne, H. Gertov, F. Sannino, and K. Tuominen, Phys. Rev. D **91**, 095021 (2015).
[17] H. Gertov, A. Meroni, E. Molinaro, and F. Sannino, Phys. Rev. D **92**, 095003 (2015).
[18] F. Feruglio, B. Gavela, K. Kanshin, P. A. N. Machado, S. Rigolin, and S. Saa, J. High Energy Phys. 06 (2016) 038.
[19] T. A. Ryttov and F. Sannino, Phys. Rev. D **78**, 115010 (2008).
[20] M. A. Luty, J. High Energy Phys. 04 (2009) 050.
[21] E. Katz, A. E. Nelson, and D. G. E. Walker, J. High Energy Phys. 08 (2005) 074.
[22] B. Gripaios, A. Pomarol, F. Riva, and J. Serra, J. High Energy Phys. 04 (2009) 070.

[23] J. Galloway, J. A. Evans, M. A. Luty, and R. A. Tacchi, J. High Energy Phys. 10 (2010) 086.
[24] J. Barnard, T. Gherghetta, and T. S. Ray, J. High Energy Phys. 02 (2014) 002.
[25] G. Ferretti and D. Karateev, J. High Energy Phys. 03 (2014) 077.
[26] G. Cacciapaglia and F. Sannino, J. High Energy Phys. 04 (2014) 111.
[27] D. Buarque Franzosi, G. Cacciapaglia, H. Cai, A. Deandrea, and M. Frandsen, J. High Energy Phys. 11 (2016) 076.
[28] R. Arthur, V. Drach, M. Hansen, A. Hietanen, C. Pica, and F. Sannino, Phys. Rev. D **94**, 094507 (2016).
[29] R. Arthur, V. Drach, A. Hietanen, C. Pica, and F. Sannino, arXiv:1607.06654.
[30] M. A. Luty and T. Okui, J. High Energy Phys. 09 (2006) 070.
[31] S. R. Coleman and E. J. Weinberg, Phys. Rev. D **7**, 1888 (1973).
[32] M. E. Peskin, Nucl. Phys. **B175**, 197 (1980).
[33] J. Preskill, Nucl. Phys. **B177**, 21 (1981).
[34] G. Aad *et al.* (ATLAS and CMS Collaborations), J. High Energy Phys. 08 (2016) 045.
[35] R. Hempfling, Phys. Lett. B **379**, 153 (1996).
[36] W.-F. Chang, J. N. Ng, and J. M. S. Wu, Phys. Rev. D **75**, 115016 (2007).
[37] C. Englert, J. Jaeckel, V. V. Khoze, and M. Spannowsky, J. High Energy Phys. 04 (2013) 060.
[38] O. Antipin, M. Mojaza, and F. Sannino, Phys. Rev. D **89**, 085015 (2014).
[39] M. Heikinheimo, A. Racioppi, M. Raidal, C. Spethmann, and K. Tuominen, Mod. Phys. Lett. A **29**, 1450077 (2014).
[40] O. Antipin, M. Redi, and A. Strumia, J. High Energy Phys. 01 (2015) 157.