Sigma-assisted natural composite Higgs

Diogo Buarque Franzosi  
Department of Physics, Chalmers University of Technology, Fysikgården, 41296 Göteborg, Sweden  
Giacomo Cacciapaglia and Aldo Deandrea  
Université de Lyon, France; Université Lyon 1, CNRS/IN2P3, UMR5822, IPNL F-69622 Villeurbanne Cedex, France

We present a simple model of dynamical electroweak symmetry breaking from a strongly interacting theory where the Higgs arises as a pseudo-Nambu-Goldstone boson. We show that the presence of a lightish scalar resonance, \( \sigma \), that mixes with the Higgs allows to relax the bounds and allows for a large misalignment angle \( s_\theta \lesssim 0.6 \), contrary to the common lore. We study the bounds in the simplest realisation with top partial compositeness, and show that the model predicts a light Higgs-like broad resonance below the TeV, which can be searched for at the LHC in large width \( \eta \eta \), \( ZZ \) final states. Finally, a light pseudo-scalar \( \eta \) is also preferred, leading to potentially observable chain decays \( h \rightarrow \eta \eta' \rightarrow ZZ\gamma\gamma \).

\[ v = f s_\theta , \]

where \( \theta \) is the misalignment angle \( \theta \). The Higgs boson \( h \) appears as a pseudo-Nambu-Goldstone boson (pNGB), thus explaining its lightness compared to the other composite states and its approximate EW doublet nature \( \theta \). In comparison, in Technicolor models \( \theta \), \( \theta \), which are matched in the limit \( \theta \rightarrow \pi/2 \), the role of the Higgs can only be played by a light singlet scalar resonance \( \theta \) or a dilaton-like light state \( \theta \).

One of the main model-building issues encountered in CH models is the fine tuning in the misalignment that is required by EW precision observables (EWPOs). Due to large corrections to the oblique \( S \) parameter \( \theta \), the compositeness scale needs to be sizeably larger than the EW scale, yielding a fairly model-independent bound \( s_\theta \lesssim 0.2 \). This, however, requires a tuning in the parameters of the model, which can happen in the top sector alone \( \theta \) or by tuning the current mass term of the underlying fermions \( \theta \). We remark that the pNGB Higgs mass is always of order \( m_h \approx f s_\theta = v \), thus its precise value is encoded in a generally incalculable strong form factor. In this letter we ask the following question: is it possible to build a feasible CH model without tuning in the misalignment potential?

Typically, it’s the top quark coupling that dominates the misalignment dynamics at low energies. If the top mass is generated by contact interactions à la extended Technicolor \( \theta \), then the natural alignment is towards the Technicolor vacuum \( s_\theta = 1 \). It has been shown in Ref. \( \theta \) that the introduction of a lightish scalar resonance \( \sigma \), that mixes with the pNGB Higgs, can alleviate the tension between EWPOs and the Technicolor vacuum. Increasing evidence of the presence of a light scalar state in theories with an infra-red conformal phase are being collected on the lattice \( \theta \) and by use of gravitational duals \( \theta \). Such state, which may or may not be a dilaton \( \theta \), necessarily mixes with the Higgs boson. Furthermore, a conformal phase, also called “walking” \( \theta \), can help alleviating the flavour issue of CH models \( \theta \) by increasing the gap between the compositeness scale and the scales of flavour violation. Finally, partial compositeness \( \theta \) has been identified as a promising mechanism to give a large mass to the top quark provided that the fermion operators that linearly mix with the top feature a large anomalous dimension in the walking window.

All the features listed above point towards realistic CH models where a light \( \sigma \) is present. Therefore, in this letter we will analyse the possibility of having a “natural” misalignment in CH models with top partners. In this perspective, it is interesting to realise that the partially composite top mass often depends on the misalignment

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1 We adopt the short-hand notation \( \sin \theta = s_\theta \), \( \cos \theta = c_\theta \) and \( \tan \theta = t_\theta \).

2 A dilaton is a pNGB associated to the spontaneous breaking of conformal invariance at quantum level.
Thus, if the potential is dominated by top loops, being proportional to $m_t^2$ the natural minimum is at $\pi/4$ ($s_\theta = 1/\sqrt{2}$) and not in the Technicolor limit. Note that this situation can be achieved if additional symmetries are imposed on the composite fermion sector, as recently proposed in Refs. [33, 39]. Besides the issue with EWPOs, a large $s_\theta$ will also induce large modification of the would-be Higgs couplings to SM states. Interestingly, such corrections are universal and model-independent: this is due to the fact that they are determined by the $\theta$-dependence of the masses. The reduced couplings to massive gauge bosons $V = W^\pm, Z$ and top (normalised to the SM values) are, therefore, equal to

$$\kappa_V = \frac{\partial V}{\partial v} = c_\theta, \quad \kappa_t = \frac{v}{f m_t} \partial_\theta m_t = c_2 \theta / c_\theta.$$  

The most serious issue is associated with the coupling to the top that vanishes at $\theta = \pi/4$. This property can be easily understood: the minimum of the potential is given by $\partial_\theta V(\theta) \propto m_t(\theta) \partial_\theta m_t(\theta)$, thus at the minimum one has either $m_t(\theta) = 0$ or $\partial_\theta m_t(\theta) = 0$. As $\kappa_t = \frac{v}{f m_t} \partial_\theta m_t$, a vacuum with non-zero top mass implies $\kappa_t = 0$. The recent detection of the $t\bar{t}h$ production channel by CMS [40] and ATLAS [41] that definitely proves $\kappa_t \neq 0$, therefore, rules out the most natural minimum. 4 The presence of a light $\sigma$ that mixes with the pNGB Higgs can alleviate both issues of EWPOs and the Higgs couplings. We will analyse this scenario in an effective field theory approach, without relying on the details of the strong sector. Nevertheless, we choose the minimal model that can be obtained in an underlying gauge-fermion theory, which is based on the symmetry breaking pattern SU(4)/Sp(4) in the strong sector [25]. The possibility to generate top partners as bound states of fermions has been analysed in Refs. [42, 43], while scalars charged under the strong dynamics were introduced in Ref. [44].

The paper is organised as follows. In sec. (I) we introduce the model and derive the relevant masses and couplings. In sec. (II) we study the constraints coming from the consistency of the theory, the Higgs couplings and EWPOs, while in sec. (III) we analyse direct searches for the heavier scalar. Finally, in sec. (IV) we offer our conclusions.

We consider a minimal model based on the SU(4)/Sp(4) coset and fermionic operators in two-index representations of SU(4) (either symmetric S or anti-symmetric A). This model can be realised in terms of a $G_{TC} = Sp(2N_c)$ gauge theory with 4 Weyl fermions $\psi = (U_L, D_L, U_L, D_L)^T$ in the fundamental representation of $G_{TC}$ [24, 45]. Additional fermions in the anti-symmetric 2-index representation of $G_{TC}$, $\chi$, carrying QCD colour and hypercharge are needed to generate top partners in the form $\psi \chi'$. 5 Note that this model is being studied on the lattice for $N_c = 2$ [47, 48]. The mass of the top comes from 4-fermion interactions of type $\psi \chi' t_{L/R}$ that, after condensation, generate a mixing of the elementary fermions with the composite ones plus additional low energy operators. The latter have been studied in detail in Refs. [33, 35].

The chiral Lagrangian describing the pNGBs — i.e. the would-be Higgs $h$, a pseudo-scalar singlet $\eta$ and the eaten NGBs $\pi^0$ — and the light singlet, is given by [25, 49]

$$\mathcal{L} = k_G(\sigma) \frac{f^2}{8} D_{\mu} \Sigma^\dagger D^{\mu} \Sigma - \frac{1}{2} (\partial_\mu \sigma)^2 - V_M(\sigma) + k_2(\sigma) \frac{y \bar{t} t_2 C}{4\pi} \Tr [\frac{(P_G^T)^\dagger P_G]}{\Sigma^\dagger} \Sigma^\dagger] + h.c. - k_1^2(\sigma) \frac{y \bar{t} t_1}{4\pi} \sigma V - k_m(\sigma) V_m,$$

where $\Sigma = v^{\Pi/2} \cdot \Sigma_0$ is the linearly transforming pNGB matrix defined around the vacuum $\Sigma_0$. The term in the second line is responsible for the top mass with the spurious $P_G$ and $P_1$ transforming in the symmetric or anti-symmetric representation of SU(4), $V_M$ is a potential for $\sigma$, and $V_{1,2,m}$ are the terms in the pNGB potential generated by the top, gauge loops and current mass respectively. Note that we do not require that $\sigma$ is a dilaton, even though our general expression can accommodate light dilaton effective lagrangians [50, 51].

We will work in a basis where none of the fields defined in eq. (3) are allowed to develop a non-zero vacuum expectation value, i.e. both pNGBs and $\sigma$ are defined around the proper vacuum. In particular, $\Sigma_0$ contains the misalignment along the Higgs direction, and the $\sigma$-potential $V_M(\sigma)$ includes a tadpole that balances up the contribution of the pNGB potential terms. Furthermore, we normalise the $\sigma$ coupling functions such that $k_i(0) = 1$ and define $k_1' = \frac{f}{m_1} |_{\sigma=0}$, $k_2' = \frac{f^2 y t_1}{m_2} |_{\sigma=0}$, etc.

The top mass operator in eq. (4), for both choices of top partner representations, gives

$$m_t = \frac{y_{LL} y_{RL} C_y f s_\theta c_\theta}{4\pi}.$$  

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3 For a survey of different top partner representations in the cosets SU(4)/Sp(4) and SU(5)/SO(5) see Refs. [35, 36] and [37].

4 The indirect probe from the top-loop induced gluon coupling is not effective as the strong sector can give additional contributions.

5 Another candidate theory is based on $G_{TC} = SO(11)$ [43], while models with scalars would have top partners in the fundamental of SU(4) [13, 16].
The three terms in the pNGB potential read \[36\]

\[
V_t = -C_t G_f^2 Y_t^2 \frac{f^2}{(4\pi)^2} \phi_0^2 + \ldots ,
\]

\[
V_g = -C_g (3g^2 + g'^2) \phi_g^2 + \ldots ,
\]

\[
V_m = -4C_m m_\psi \phi_0 + \ldots ,
\]

where \(m_\psi\) is a common mass for the two underlying fermions \[23\] and we only show the field-independent terms (a more complete expansion can be found in Ref. \[36\]). For later convenience we define the gauge form factor in terms of the top one as

\[
\delta \equiv C_g (3g^2 + g'^2) \frac{2}{2C_t'} , \quad C_t' \equiv C_t G_f^2 Y_t^2 \frac{f^2}{(4\pi)^2} .
\]

This leads to the minimum condition

\[
\frac{\partial V}{\partial \theta} = 2f^3 s_\theta (f C_t' (c_{2\theta} - \delta) c_\theta - 2C_m m_\psi) = 0 .
\]

The zero at \(s_\theta = 0\) would imply that the EW symmetry is unbroken (\(\theta = 0\)). However, as we expect the top loop to drive the potential to break the EW symmetry (i.e., \(C_t > 0\)), the minimum of the potential should sit at

\[
C_m m_\psi = \frac{1}{2} C_t' f c_\theta (c_{2\theta} - \delta) .
\]

Note that the gauge loops typically do not break the gauge symmetry itself, thus \(\delta > 0\), and the existence of such solution is fully encoded in the contribution of the mass, which cannot be too large. Also, for \(m_\psi = 0\), we find \(c_{2\theta} = \delta\). Assuming that the top interactions dominate, i.e. \(\delta < 1\), the most likely minimum is at \(c_{2\theta} \approx 1\), which is therefore the most “natural” misalignment in this class of models. Moving away from it would require either to enhance \(\delta\) by suppressing the top contribution to the potential, or by adding a sizeable mass \(m_\psi\), thus incarnating the fine tuning issue of CH models.

**A. Masses**

The would-be pNGB Higgs \(h\) mixes with \(\sigma\): the mass matrix is given by

\[
- \frac{1}{2} (h \quad \sigma) \begin{pmatrix} m_h^2 & \frac{1}{2} (A m_h^2 + B \tilde{m}_h^2) \\ \frac{1}{2} (A m_h^2 + B \tilde{m}_h^2) & m_\sigma^2 \end{pmatrix} (h \quad \sigma) ,
\]

where

\[
m_h^2 = 2 C_t s_\theta \frac{\delta (\delta - 3c_{2\theta} - 2)}{c_\theta^2} ,
\]

\[
m_\sigma^2 = -8 u C_m m_\psi \equiv m_\eta^2 c_\theta^2 \frac{\delta (\delta - c_{2\theta})}{s_\theta^2} ,
\]

\[
A = \frac{c_{2\theta}}{2s_\theta} (k'_G - 2k'_t) ,
\]

\[
B = \frac{s_\theta}{4} [2(k'_m + k'_t) - 3k'_G] + \frac{t_\theta}{4} (k'_G - 2k'_t) ,
\]

while we consider the mass of \(\sigma\) as a free parameter as it receives a potentially dominant contribution from the potential, \(m_\sigma^2 \sim V_h' / f^2 + O(m_h^4)\). We have also imposed the minimum condition in eq. \[10\] by replacing \(m_\psi\) in all expressions. Physically the term \(m_\eta^2\) corresponds to the contribution of \(m_\psi\) to the mass of the singlet \(\eta\). For the symmetric \(S\) spurion, this is the only mass term, while for the anti-symmetric \(A\) an additional term is generated:

\[
m_\eta^2 = m_\eta^2 + \frac{\epsilon_A}{4} \left( m_\eta^2 - \tilde{m}_\eta^2 \right) ,
\]

where \(\epsilon_A \leq 1\) for the anti-symmetric top spurion, and zero otherwise. The actual value of \(\epsilon_A\) depends on the embedding of the top singlet \(t_R\).

The masses of the physical eigenstates \(h_{1,2}\) are given by

\[
m_{h_{1,2}}^2 = \frac{1}{2} \left\{ m_h^2 + m_\eta^2 \pm \sqrt{(m_h^2 - m_\eta^2)^2 + 4(A m_h^2 + B \tilde{m}_h^2)^2} \right\} .
\]

In principle, either state can play the role of the 125 GeV Higgs, however we expect the \(\sigma\) to be heavier than the EW scale and its couplings to be far from the SM ones. Thus, we will conservatively associate the Higgs boson with the lighter state, \(m_{h_1} = 125\) GeV, and make sure that in the limit of no mixing, where \(A, B \to 0\), \(m_{h_1} \to m_h\) and \(m_\sigma \to m_{h_2}\). This is justified in the context of a composite heavier sector. The relation between the mass eigenvalues and the mass parameters in the mixing matrix of eq. \[12\] can thus be inverted as follows:

\[
m_{h_{1,2}}^2 = \frac{1}{2A^2 + 1} \left\{ m_{h_1}^2 + m_{h_2}^2 - 2AB \tilde{m}_h^2 \pm \sqrt{(m_{h_1}^2 + m_{h_2}^2 - 2AB \tilde{m}_h^2)^2 + 4(A^2 + 1)(B^2 \tilde{m}_h^2 + m_{h_1}^2 m_{h_2}^2)} \right\} .
\]

We remark that both solutions are real and positive, provided that the argument of the squared root is positive.
The potential also generates trilinear couplings among the pNGBs:
\[
g_{h^3} = \frac{3m_h^2 c_2 v}{c_\theta} \frac{c_2 - 2}{c_\theta} + \frac{m_h^2 s_2^2}{c_\theta},
\]
\[
g_{h\eta^2} = \frac{m_h^2 s_2^2}{c_\theta} + \frac{m_h^2 s_2^2 (c_2 + 7)}{6v c_\theta}.
\]

The couplings of \( \sigma \) to two pNGBs are
\[
g_{\sigma h^2} = k_G' \frac{p_1^* p_2}{v} s_\eta + \mathcal{O} \left( \frac{m_\eta^2}{v^2} \right),
\]
\[
g_{\sigma \eta^2} = k_G' \frac{p_1^* p_2}{v} s_\eta + \mathcal{O} \left( \frac{m_\eta^2}{v^2} \right).
\]

where the first term in both comes from the kinetic term and corrections come from the pNGB potential. The momenta \( p_l \) refer to the pNGBs, and we recall that on shell \( p_1 p_2 \approx m_\eta^2/2 \). This contribution ensures that in the large mass limit the equivalence theorem is attained and the partial width of \( \sigma \) into the 5 pNGBs are the same.

For heavy masses \( m_{h_2} \approx m_\sigma \gg v \) the total width of the heavier scalar \( h_2 \) is dominated by the \( \sigma \) component, and is equal to
\[
\Gamma/m_{h_2} \approx \frac{5k_G'^2}{128\pi f^2} m_{h_2}^2 + \frac{k_i^2 N_c m_{h_2}^2}{8\pi f^2}.
\]

The first term can be recognised as the partial width into pNGBs, while the second is the partial width in tops. For our perturbative treatment of the mixing between \( h \) and \( \sigma \), we need to make sure that the width remains small or comparable to the mass of the scalar. This will be one of the constraints we consider in the following section.

### III. Constraints

If we only focus on the scalar sector, \( h_{1,2} \), the model has 6 free parameters: the mass of the heavier scalar, \( m_{h_2} \), the mixing parameters \( A \) and \( B\tilde{m}_\eta^2 \) (cf. eq. (12)), \( \delta \) (cf. eq. (9)), the misalignment angle \( \theta \) and the coupling \( k_G' \). We recall that we have used the mass \( m_{h_2} \) to ensure the minimum condition, and traded the two \( \sigma \)-couplings \( k_i' \) and \( k_i' \) with \( \delta_A \) and \( B \). Finally, \( \tilde{m}_\eta^2 \) is related to the mass of the singlet pNGB \( \eta \) (cf. eq. (17)).

In this section and the following one we will probe the parameter space of the model by imposing constraints deriving from various considerations:

i - Consistency of the theory, which includes perturbativity of the couplings and perturbative unitarity of pNGB scattering;

ii - Higgs property measurements, namely its couplings and total width;

iii - EWPOs;
iv - Direct search of the heavy scalar.

The main goal of this exercise is to determine how large \( \theta \) is still allowed to be once all constraints are taken into account.

A. Consistency conditions

We consider 3 consistency conditions on the parameter space of our effective model.

Firstly, we require that all couplings are perturbative. To this effect, we demand that all the \( \sigma \)-couplings respect \(|k'_i| < 4\pi, i = t, G, m\).

Secondly, we demand that the pNGB scattering remains perturbatively unitary up to the condensation scale \( 4\pi f \). This requirement is connected to the one above, as we expect that the \( \sigma \) plays a crucial role in taming the growth with energy of the amplitude, like in QCD \[52\]. We will base our estimate on the leading order calculation, even though radiative corrections typically tend to increase the amplitude and push the resonance mass to lower values \[53\]. Neglecting effects from the potential, which are irrelevant at high energies, the asymptotic behaviour of the pNGB scattering in the sigma channel (projection on zero isospin and angular momentum, \( I = J = 0 \)) is given by

\[
a^{(0)}_{A0}(s) \approx \frac{s}{16\pi f^2}.
\] (32)

We thus require that the mass of the \( \sigma \) lies below the energy scale where the above amplitude grows larger than 1. This allows us to define a parameter, \( \gamma \), that measures how far the mass is from the boundary of perturbative unitarity loss

\[
\gamma \equiv \frac{m_{h_2}}{4\sqrt{\pi}f} \lesssim 1.
\] (33)

In the following we will fix \( \gamma \) and replace the value of the heavy scalar mass. Note that the annihilation in \( t\bar{t} \) may also be relevant at low scales, thus pushing the unitarity loss scale lower.

Thirdly, we require that the heavy scalar width remains small compared to the mass. A width for \( h_2 \) much larger than the mass would transform the state into a broad excess instead of a resonance, while jeopardising the perturbative description of the mixing. Although a broad state is a plausible scenario, we require \( \Gamma/M < 1 \) in order to trust our calculations for the Higgs \( h_1 \).

The regions with non-perturbative \( \sigma \)-couplings are shown in blue in Fig. 2 for fixed values of some parameters (specified in the caption). The regions with vertical hash lines correspond to \( k'_{1\nu} \), while the region with diagonal hash lines corresponds to \( k'_{t} \), at the bottom of the plots. The black lines mark contours of fixed \( h_2 \) width \( \Gamma_{h_2}/m_{h_2} = 0.3, 0.5, 1 \), with the smaller value on the internal contour. Perturbative unitarity is ensured by fixing \( \gamma = 0.2 \) (left and right panels) or \( \gamma = 0.4 \) (middle panel).

B. Higgs measurements

Since the discovery of the Higgs boson, both ATLAS and CMS have been measuring its couplings with increasing precision. These measurements provide relevant limits on any model that modifies the Higgs sector. We already discussed in the previous section how the couplings of \( h_1 \) to \( V = W^\pm, Z \) and the top are modified. In general, also the couplings to light fermions, like the bottom and tau, will be affected, while direct contributions of the strong sector may affect the couplings to gluons and photons, which are loop induced in the SM. However, the details are model dependent. Here we want to be conservative, so we will extract only bounds on \( \kappa_V \) and \( \kappa_t \) that are independent on other measurements. For the coupling to vectors, we use the combined fit of ATLAS and CMS after Run-I \[54\] and extract the bound on \( \kappa_V \) from the most general fit. The coupling to tops is bound indirectly by the measurement of the gluon fusion cross section. However, if we allow for a generic contribution to the gluon coupling from new physics, the only solid bound comes from the observation of the \( t\bar{t}h \) production mode \[40,41\]. We thus translate the most stringent bound on the signal strength, \( \mu = 1.26^{+0.31}_{-0.26} \) from CMS, to a bound on \( \kappa_t \). The Higgs coupling bound we impose are, therefore,

\[
\kappa_V = 1.035 \pm 0.095, \quad \kappa_t = 1.12^{+0.14}_{-0.12}. \tag{34, 35}
\]

Note that Run-II bounds on the couplings to vectors are becoming more constraining, however a combination is still not available and we refrain to do it as taking into account systematics of the experiments cannot be reliably done. The regions excluded at 2 sigma level (95% C.L.) are shown in red in Fig. 2, with the vertical hash lines for the top couplings and diagonal ones for the vectors. The plot thus highlights the importance of the measurement of the \( t\bar{t}h \) production for CH models. The constraints, especially coming from \( \kappa_t \), are more severe for \( \delta_A \) far from 1, for instance the left plot would have the whole EWPO valley between the green shaded area covered if \( \delta_A = 0.7 \).

The Higgs properties are also affected by the presence of the pseudo-scalar singlet \( \eta \), whose mass can be below the threshold of the decay \( h \to \eta \eta \). The global fit of Ref. \[54\] provides the following bound on the branching ratio of the Higgs into undetected non-SM states

\[
B_{BSM} < 0.32, \tag{36}
\]

at 2 sigmas. This imposes a strong constraint on the parameter space due to the Higgs decay \( h \to \eta \eta \). We estimate the Higgs total width as

\[
\Gamma_h = (c_{\beta\alpha})^2(\Gamma_{\beta} + \Gamma_{\gamma}) + (\kappa_V^{h})^2\Gamma_V + \Gamma_{\gamma} + \Gamma_{BSM}, \tag{37}
\]

where \( \Gamma_{\gamma} \) are the SM partial widths. This expression assumes that the bottom and tau get their masses from
a bilinear term, as in Ref. [25]. For the decay into gluons we use the SM value as a first approximation. The bound from eq. (36) is shown as a dashed line in the left panel of fig. 2 for \( m_\eta = m_\eta = 0 \). We see that it allows to exclude the large values of \( \theta \), thus it is necessary to give mass above threshold to the singlet \( m_\eta > m_h/2 \). In the anti-symmetric case, this condition might be fulfilled even for vanishing underlying fermion mass, according to eq. (17). Note that off-shell decays \( h \to \eta \eta^* \) might also be sizeable and give interesting final states as \( \eta \to Z \gamma \) below the WW threshold [22]. Thus the final state \( h \to ZZ \gamma \gamma \) may be a smoking gun for this model.

\[ \Delta S = \frac{1 - (\kappa_V^{h_1})^2}{6\pi} \log \frac{\Lambda}{m_{h_1}} - \frac{(\kappa_V^{h_2})^2}{6\pi} \log \frac{\Lambda}{m_{h_2}} + \Delta S_{\text{TC}}, \]

where \( \Delta S_{\text{TC}} = N_D s_3^2/(6\pi) \) is the contribution of the strong sector [18, 55], with the \( N_D \) factor counting the number of EW doublets in the underlying theory \( N_D = 2N_c \) in the minimal model introduced in sec. (II) and we use \( \Lambda = 2\pi^2 f^2 \) as the compositeness scale.

Vector and axial-vector resonances are also known to contribute to the oblique parameters and cause cancellations [56, 57]. The contribution in the model under study was computed in Ref. [58] and amounts to replacing the term \( \Delta S_{\text{TC}} \) with

\[ \Delta S_{\text{TC}} \to \Delta S_R = \frac{16\pi (1 - r^2) s_3^2}{2(g^2 + g^2) - g^2(1 - r^2)s_3^2} \bar{\kappa}_G \bar{\kappa}_h \bar{\kappa}_{h} \].

with \( \bar{g} \) and \( r \) being non-perturbative parameters of the chiral Lagrangian. A cancellation thus happens for \( r > 1 \), where the new correction to \( S \) is negative. In fig. 2 we show in green the excluded region at 95% CL for \( r = 1.1 \) and \( \bar{g} = 3 \), where a large region is allowed for \( k_G' \) of order unity, both positive and negative. \(^6\) Unitarity and dispersion relation arguments can also be applied to the vector resonances and, interestingly, they indicate values \( |a_\rho| \approx 1 \), where \( a_\rho \equiv \sqrt{2M_f/(1 - r^2)} \) [53], and thus \( r \approx \pm 0.2 \) or \( r \approx \pm 0.9 \) for \( \bar{g} = 3 \). For larger values of \( r \) the partial widths of the vectors into longitudinal bosons rapidly grow and leads to excessively broad resonances, \( \Gamma_\rho/M_\rho = 1 \) for \( r \approx \pm 0.1 \) or \( r \approx 0.80 \). For low values of \( r \) the decay into longitudinal bosons are uncomfortably suppressed to be motivated in a strongly interacting theory. We therefore regard \( r = 1.1 \) as a well-motivated value based on the previous considerations.

The EWPOs are clearly very constraining, however we remark that the precise excluded region depends crucially on the details of the strong dynamics, as indicated by the effect of the vector resonances. Other resonances, like the top partners, may also contribute, thus further changing this region and, potentially, opening up a wider parameter space.

\[ \Delta T = \frac{-3(1 - (\kappa_V^{h_1})^2)}{8\pi c_w^2} \log \frac{\Lambda}{m_{h_1}} + \frac{3(\kappa_V^{h_2})^2}{8\pi c_w^2} \log \frac{\Lambda}{m_{h_2}}, \]

IV. DIRECT SEARCHES FOR THE HEAVY SCALAR AT THE LHC

The main signature of this model is the presence of a second “Higgs” \( h_2 \) that can be observed at the LHC. Its mass is a free parameter, however, as we have seen, its mass is limited by requiring perturbative control of the effective theory. The production mechanisms are the same as for the SM Higgs, namely gluon fusion (ggF) and vector boson fusion (VBF), with associated production with tops to a lesser extent.

The production of \( h_2 \) via ggF is difficult to estimate due to its loop-nature: besides the top loop generated by the coupling in eq. (20), the strong dynamics can generate an additional direct coupling to gluons. To understand the structure of the latter, we analyse the possibility that it is generated dominantly by a loop of a heavy top-like resonance \( T \), aka top partner. The coupling of the \( \sigma \) will have the form

\[ k_T(\sigma) M_T T \bar{T} = \left( M_T + k_T f T \sigma + \ldots \right) T \bar{T}, \]

where the top partner mass \( M_T = g_T f_T \), with \( g_T = O(1) < 4\pi \). Thus, one can define a reduced coupling

\[ \kappa_T^2 = k_T g_T s_3, \]

which is explicitly suppressed by a power of the misalignment angle. This suppression will appear in the effective coupling to gluons.

To estimate the production cross section of \( h_2 \) we use the N^3LO result for ggF [60] and the NNLO for VBF.
production \cite{61, 62}, and rescale the SM Higgs production cross section as follows:

$$
\sigma = \sigma_0^H \left| \frac{h^2_{t} A_{t} (\tau_3) + h^2_{g} }{A_F (\tau_3)} \right|^2 + \sigma_0^{VBF} (\kappa^2_{h^2})^2 ,
$$

where $A_{t} (\tau_3)$ is the standard loop amplitude for the top quark in the SM. Following the argument above, in the following we will fix the strong dynamics contribution to the coupling to gluon as $\kappa^2_{h^2}/\delta_0 = $ const.

The strongest experimental constraint on a heavy Higgs-like resonance comes from ZZ searches. The main issue with reinterpretation of the experimental result is due to the fact that $h_2$ in this model tends to be very broad. In Ref. \cite{57} the CMS collaboration considered broad scalar resonances decaying into ZZ final states, with widths up to $\Gamma/m < 0.3$. We can thus extract expected exclusions by simply comparing the production rates of the ZZ final states directly with the experimental results. This is shown in fig. \text{(3)} for the parameters specified in the caption. The yellow region is thus excluded at 95\% CL, and we show in dashed the result for a narrow resonance and in solid for $\Gamma/m = 0.3$. The closeness of the two curves proves that the effect of the width is not very important. In the right side of each plot we show the mass of $h_2$, which is not fixed in the plot but related to $c_{2g}$ by eq. \text{(41)} once we fix $\gamma = 0.2$ or 0.4. The increase of the mass for $\theta \rightarrow 0$ is compensated by an increase of the branching ratio into gauge bosons in the same limit, thus explaining why we do not lose too much sensitivity for larger $h_2$ masses.

In this scenario, decays into tops are also relevant: for instance, in the allowed region of the left panel of fig. \text{(3)}, the branching ratio of $h_2 \rightarrow t\bar{t}$ lies between 70\% and 90\%. However, this is a very challenging search due to large interference between signal and background. We have used the framework developed in Ref. \cite{61} to access the power to discover the heavy scalar via top pair production taking these specific parameters as benchmark. The analysis is based on the comparison of the measurement of the differential cross section of the $t\bar{t}$ invariant mass distribution in $t\bar{t}$ process at particle level done by the ATLAS collaboration in a resolved \cite{65} and a boosted regime \cite{66}. Only for large values of $\kappa_{h^2}^2 \gtrsim 4\sigma_0$ this search becomes competitive with the ZZ search. In the right panel of fig. \text{(3)} the bound from top pair production appears as we chose $\kappa_{h^2}^2 = 5\sigma_0$. The exclusion is derived by a line-shape analysis on the $m(t\bar{t})$ distribution, which assumes that the data fit exactly the SM prediction for collisions at a centre of mass energy of $\sqrt{s} = 13$ TeV and integrated luminosity of 20 fb$^{-1}$. We used a $m(t\bar{t})$ resolution of 40 GeV, uncorrelated systematic errors of 15\% on all bins and a theoretical uncertainty of 5\%. However, a dedicated experimental analysis searching for this kind of broad resonances, with variable values of the total width and effective gluon couplings, would be necessary to ascertain the reach at the LHC.

V. DISCUSSION AND CONCLUSIONS

We have analysed the constraints on a simple scenario of composite Higgs with top partial compositeness. The presence of a light scalar resonance that mixes with the would-be Higgs allows to release the strong bounds on the misalignment angles. In fig. \text{(2)} we show, in three benchmark points, that it is still possible to allow large values of the misalignment angle down to $c_{2g} \approx 0.2$, i.e. $s_0 \approx 0.6$, for masses of the second Higgs around 600 GeV. The indirect bounds mostly depend on the value of the mass of the heavy scalar. By increasing it, as shown in the middle plot, the constraints from perturbativity of the couplings tend to exclude large misalignment angles. A minor dependence arises on the mass of the pseudo-scalar singlet $\eta$, as shown comparing the left and right plots. However, for light $\eta$, decays $h \rightarrow \eta \eta$ tend to rule out most of the parameter space. Conversely, above threshold, measurable rates $h \rightarrow \eta \eta^*$ might occur, with final state $h \rightarrow ZZ^* \gamma \gamma$ deriving from the anomalous couplings of $\eta$. This could...
be one of the smoking gun signatures of this model.

In order for this mechanism to work, a second Higgs-like state is predicted with a mass below 1 TeV, which can thus be produced at the LHC. Searches for this state are complicated by the fact that the coupling to gluons cannot be predicted in an effective field theory framework and by the presence of a sizeable width. We tested the reach of ZZ resonant searches at large width ($\Gamma/m < 30\%$), which can potentially test the large misalignment angle region. Searches for $t\bar{t}$ final states, complicated by the large interference with the SM background, are also relevant for large couplings to gluons.

The need for a small misalignment angle in composite pseudo-Nambu-Goldstone Higgs models is often used as a measure of the fine tuning in the models. We have shown that the angle can be larger than previously thought provided the presence of a heavier scalar state. In the minimal scenario we studied, the misalignment is determined by the interplay between the top and gauge loops and a current mass for the underlying fermions. The latter also gives mass to the pseudo-scalar singlet $\eta$. The ratio $\delta$ between the gauge and top loops play a crucial role in determining the misalignment angle. In fig. 2 we show lines of fixed $\delta$ in blue: they mainly depend on $\tilde{m}_\eta$, cf. eq. (17), and only in a minor way on the other parameters via the value of $m_t$, cf. fig. 1. We see that increasing the value of the current mass, i.e. $\tilde{m}_\eta$, increases the sensitivity of $\theta$ on the value of $\delta$, thus increasing the necessary tuning in the potential. In particular, as shown in the right panel of the figure, too large a mass requires negative $\delta$ in order to achieve large $\theta$. However, $\delta$ is expected to be positive from arguments based on the stability of the potential. These facts seem to point towards light $\eta$, which can thus give visible signals via off-shell Higgs decays, as discussed above.

In conclusion, we have shown that large values of the misalignment angle in composite Higgs models are not yet excluded, and that they predict the presence of a light Higgs-like scalar broad resonance that can be searched for at the LHC in ZZ and $t\bar{t}$ final states. A preference for a light pseudo-scalar in the minimal scenario may also provide smoking gun signatures $h \to ZZ^*\gamma\gamma$ via on- and off-shell cascade decays $h \to \eta\eta'$, $\eta \to Z\gamma$.

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