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Light ’t Hooft top partners

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Vectorlike quarks, usually dubbed top partners, are a common presence in composite Higgs models. Being composite objects, their mass is expected to be of the order of their inverse size, that is the condensation scale of the new strong interactions. Light top partners, while not being a generic prediction, are, however, often considered in phenomenological models. We suggest that their lightness may be due to the matching of global ’t Hooft anomalies of the underlying theory. We check this mechanism in explicit models, showing that, in one case, composite fermions with the quantum numbers of the top quark obtain a mass which is controlled by a soft breaking term and can be made parametrically small.

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I. INTRODUCTION

Compositeness is an attractive hypothesis to solve the Standard Model (SM) hierarchy problem. In modern composite Higgs models (CHMs), the Brout-Englert-Higgs doublet arises as a pseudo Nambu-Goldstone boson (pNGB) due to a new strongly interacting sector that breaks spontaneously a global symmetry, not shared by the SM fields, to a properly chosen subgroup. Since it is a Goldstone, shift symmetry forces the Higgs potential to depend on explicit sources of the symmetry breaking and to vanish if the latter are sent to zero, namely if the new physics sector decouples from the SM. Partial compositeness [1] is typically a key ingredient of pNGB Higgs models to transmit the electroweak symmetry breaking from the strong sector to the elementary sector, namely to the SM fermions, or at the very least to the top quark. For reviews see Refs. [2–4] and references therein. Partial compositeness implies a linear mixing of each chirality of the SM fermions with an operator of the strong sector with matching quantum numbers; this usually implies the presence of vectorlike fermionic partners belonging to the realm of the resonances, i.e. bound states of the strongly coupled sector. Since partial compositeness is effectively a seesaw mechanism [5], large SM masses are favored for large ratios of mixing over masses; this implies that the resonance coupling to the top needs to be light, if we want to keep mixing couplings perturbative and small perturbations of the strong sector. Moreover, these resonances contribute to the Higgs potential via the linear mixings and, under some assumptions [6], help tame its sensitivity to high energy scale physics [7]. Light top partners are generically favored by considerations on the Higgs potential [6,12,13], although there are ways to evade this conclusion [14,15]. Light uncolored fermions are also needed in some composite twin Higgs models [16,17]. Finally the possibility of light top partners is interesting because it can be directly tested at the LHC; Run I established limits around 800 GeV on their masses [18–20], while Run II will explore masses up to 1.4 TeV [21,22] (or at most up to 2 TeV [23]), and heavier fermions are certainly outside the reach of the machine.

The masses of fermionic resonances in a generic strongly coupled theory are expected to be at the same scale as the mass of vector resonances, i.e. in the multi-TeV range. For scalar resonances, appearing as pNGBs, a shift symmetry can bear the responsibility of their small mass; however, no such symmetry is present for fermions. There exist extra dimensional constructions where this single scale degeneracy is relieved and Kaluza-Klein fermionic states are lighter than others, as in Refs. [24–28]. Following the AdS/CFT correspondence, such theories are dual to four dimensional (4D) strongly coupled conformal field theories and the light states interpreted as zero modes of spin-$1/2$ operators [29] (see also Ref. [30] for a supersymmetric analysis); the presence of such states crucially depends on the dimension of the operator in the conformal theory. In this paper we want to focus on the possibility that the top partners emerge as bound states of an underlying 4D confining dynamics (not necessarily conformal), which only contains matter fermions; in such a case, it is nontrivial to obtain feasible operators [31]. A simple way out that allows one to protect fermion masses could be to introduce supersymmetry [32–34], at the price of reintroducing fundamental scalars.[1]

In this paper we want to stick to classes of models with purely fermionic components and propose the possibility that light composite fermions may be present in the spectrum due to the ’t Hooft anomaly matching [36].

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[1] A model of composite vectorlike fermions, formed by a fermion and a scalar, which become light for the binding coupling close to a critical value is studied in [35].

This possibility, and related ideas, was widely employed in the 1980s [37,38] in the quest for light composite fermions that may play the role of the SM fermions; examples can be found in Refs. [39–42]. The template model we will consider is based on a strong sector based on a hypercolor (HC) gauge group G_{HC} with one, or more, species of fermions transforming under different representations of the HC group. The global symmetry $\mathcal{G}$ of the model is thus determined by the number of species and their multiplicity. We also assume that the confinement of the HC group generates one or more fermion condensates which are nonvanishing in the vacuum of the theory, thus breaking $\mathcal{G}$ to a subgroup $\mathcal{H}$. The 't Hooft anomaly matching is based on the fact that $\mathcal{H}$ may suffer from global anomalies in the underlying theory, which depend on the details of the underlying dynamics. In the confined phase, the value of the anomaly should be matched by the presence of massless composite fermions, transforming under suitable representations of $\mathcal{H}$. The matching is highly nontrivial as the representations of $\mathcal{H}$ in the spectrum are constrained by the structure of the fermionic content in the underlying theory, while the anomaly in the unconfined phase depends on the number of hypercolors. If on the other hand no solution is found, the breaking pattern $\mathcal{G} \to \mathcal{H}$ cannot be realized, and additional condensates must be turned on. Further conditions may apply if one considers the decoupling limit of the underlying fermions; however, this relies on the absence of phase transitions, and, following Ref. [41], we will not consider it here. Once the low energy theory contains 't Hooft composite fermions, a mass term for them can then be generated by adding an explicit breaking of $\mathcal{H}$, for instance in the form of a mass term m_χ for the fundamental fermions; the top partner masses must therefore vanish for vanishing m_χ, and their value can be made parametrically smaller than the mass of other resonances. Hence, such a situation generates technically natural light masses.

We should stress, however, that finding a solution for the 't Hooft matching does not imply that such spectrum is realized; the vacuum is dynamically chosen by the strong sector, and it can only be determined by nonperturbative techniques, like on the lattice, or in suitable limits of the theory, such as large N_{HC} expansion. Furthermore, the presence of an explicit, albeit small, breaking of the global symmetry may destabilize the vacuum.

In the following we illustrate the mechanism for a few realistic models: in a case where it works in Sec. II and in two models where it is not applicable in Sec. III, while in Sec. IV we conclude. We leave for a future work a more systematic study, characterizing theories in terms of the possibility of matching global anomalies.

II. MINIMAL MODEL SU(N_Q) → SP(N_Q)

The first model under consideration consists of a $G_{HC} = \mathrm{Sp}(2N_c)$ gauge theory with two species of chiral fermions: Q in the fundamental and χ in the two-index antisymmetric, where the absence of the Witten anomaly requires even N_Q. The largest global symmetry is therefore $\mathcal{G} = \mathrm{SU}(N_Q) \times \mathrm{SU}(N_\chi) \times \mathrm{U}(1)$. The model is summarized in Table I, and the minimal case of CHMs corresponds to $N_Q = 4$ and $N_\chi = 6$ [43].

TABLE I. Fermionic field content of the first model.

	$\mathrm{Sp}(2N_c)$	$\mathrm{SU}(N_Q)$	$\mathrm{SU}(N_\chi)$	$\mathrm{U}(1)$
Q	$\square$	N_Q	1	1
χ	$\begin{smallmatrix}\square\\\square\end{smallmatrix}$	1	N_χ	$-\frac{N_Q}{2N_\chi(N_c-1)}$

In the confined phase, there exist potentially two fermion condensates that spontaneously break the global symmetries when assuming a nonzero value on the vacuum: $\langle QQ\rangle$ breaks $\mathrm{SU}(N_Q) \times \mathrm{U}(1) \to \mathrm{Sp}(N_Q)$, thus leading to a pNGB Higgs in the spectrum and $\langle\chi\chi\rangle$ would break $\mathrm{SU}(N_\chi) \times \mathrm{U}(1) \to \mathrm{SO}(N_\chi)$, where the QCD color $\mathrm{SU}(3)_c \subset \mathrm{SO}(N_\chi)$. In the following we will consider two phases that lead to a potentially interesting pNGB Higgs model.

A. Phase $\langle QQ\rangle \neq 0$, $\langle\chi\chi\rangle \neq 0$

In this case, the unbroken global symmetry is $\mathcal{H} = \mathrm{Sp}(N_Q) \times \mathrm{SO}(N_\chi)$, which has no global anomalies. Therefore, all composite fermions are heavy, and there is no symmetry reason why some of them should be parametrically lighter than the other resonances. On the other hand, the spontaneous breaking of $\mathrm{SU}(N_\chi)$ implies the presence of light colored scalars [44].

B. Phase $\langle QQ\rangle \neq 0$, $\langle\chi\chi\rangle = 0$

In this phase, $\mathcal{H} = \mathrm{Sp}(N_Q) \times \mathrm{SU}(N_\chi)$. As the global symmetry containing the QCD color is unbroken, no light colored pNGBs will be present in the spectrum. In the underlying theory, $\mathrm{SU}(N_\chi)$ has a global anomaly proportional to

$$A_{\mathrm{SU}(N_\chi)^3} = \dim(\begin{smallmatrix}\square\\\square\end{smallmatrix}) = (2N_c+1)(N_c-1)\,. \tag{1}$$

This anomaly should be matched by the composite fermions in the confined phase; considering the lowest dimensional operators, we have three-fermion states $QQ\chi$ ($\bar{Q}\,\bar{Q}\chi$) and $Q\bar{Q}\bar{\chi}$, where all states have the same chirality. For $N_c > 2$, bound states $\chi\chi\chi$ and $\chi\bar{\chi}\bar{\chi}$ are also possible; however, they do not couple to the Higgs in the $\mathrm{SU}(N_Q)/\mathrm{Sp}(N_Q)$ coset and are therefore not suitable to be top partners. Bound states with a larger number of fermionic components are also possible; however, we will not consider them here as they are likely to quickly decay into lower dimensional states and may thus be highly unstable. In this model, all top partners transform as either the fundamental (**F**) or antifundamental ($\bar{\mathbf{F}}$) of $\mathrm{SU}(N_\chi)$, and

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TABLE II. Three fermions bound states of the model with group properties with respect to the global flavor group and the unbroken subgroups.

	$SU(N_Q)\times SU(N_\chi)$	$Sp(N_Q)\times SU(N_\chi)$	$d_{Sp(N_Q)}$
χQQ	$(\mathbf{A},\mathbf{F})$	$(\mathbf{1},\mathbf{F})$	1
		$(\mathbf{A},\mathbf{F})$	$\frac{N_Q(N_Q-1)}{2}-1$
$\chi\bar{Q}\,\bar{Q}$	$(\mathbf{S},\mathbf{F})$	$(\mathbf{S},\mathbf{F})$	$\frac{N_Q(N_Q+1)}{2}$
$\bar{\chi}\,\bar{Q}\,Q$	$(\mathbf{1},\bar{\mathbf{F}})$	$(\mathbf{1},\bar{\mathbf{F}})$	1
	$(\mathbf{Adj},\bar{\mathbf{F}})$	$(\mathbf{A},\bar{\mathbf{F}})$	$\frac{N_Q(N_Q-1)}{2}-1$
		$(\mathbf{S},\bar{\mathbf{F}})$	$\frac{N_Q(N_Q+1)}{2}$

thus their contribution to the $SU(N_\chi)^3$ anomaly is simply given by the multiplicity of the representation under $Sp(N_Q)$; as Q and $\bar{Q}$ transform as the fundamental of $Sp(N_Q)$, all bound states will contain a singlet $\mathbf{1}$, a two-index symmetric $\mathbf{S}$, and a two-index antisymmetric $\mathbf{A}$, as shown in Table II. The anomaly matching condition can thus be simply expressed as

$$n_{\mathbf{1}}+\left(\frac{N_Q(N_Q-1)}{2}-1\right)n_{\mathbf{A}}+\frac{N_Q(N_Q+1)}{2}n_{\mathbf{S}} = (2N_c+1)(N_c-1), \tag{2}$$

which gives a nontrivial relation between the number of flavors of Q and number of hypercolors. In the above equation, $n_{\mathbf{X}}$ is the difference of the number of fundamental and antifundamentals of $SU(N_\chi)$ in the representation $\mathbf{X}$ of $Sp(N_Q)$. Due to the presence of singlets of $Sp(N_Q)$, the above condition always has a trivial solution when the multiplicity of hypercolors is matched by the number of massless singlets.

Very attractive solutions can be achieved if a relation between the number of Q-flavors and the number of hypercolors is present; for instance, a single antisymmetric of $Sp(N_Q)$ is sufficient if $N_Q=2N_c$. In the minimal model with $G_{HC}=Sp(4)$, this singles out top partners in the $\mathbf{A}=\mathbf{5}$ of $Sp(4)\simeq SO(5)$. Interestingly, there are no solutions where the only massless fermion is an $\mathbf{S}$ nor when only an $\mathbf{A}$ and a singlet are present; these cases would correspond to a single complete $SU(N_Q)$ representation. In the minimal model, we also found a simple solution containing one $\mathbf{S}=\mathbf{10}$ in the fundamental of SU(6) and one $\mathbf{A}=\mathbf{5}$ in the antifundamental, corresponding to $n_{\mathbf{S}}=1=-n_{\mathbf{A}}$ and $n_{\mathbf{1}}=0$.

The global $SU(N_\chi)$ can be explicitly broken to $SO(N_\chi)$ by giving a gauge invariant mass to the χ's, so that the massless fermions will acquire a mass that scales with m_χ. Note also that the global $Sp(N_Q)$ potentially suffers from a global Witten anomaly [45]; however, it identically vanishes in this model.

III. OTHER MODELS

In this section we discuss other models that recently appeared in the literature; in particular we check whether this mechanism to protect the mass of top partners can work in the models considered in Refs. [46] and [47]. Unfortunately this is not the case, meaning that 't Hooft anomalies cannot be matched in one case and identically vanish in the other.

The model in Ref. [46] consists of a $G_{HC}=SU(4)$ gauge theory and three species of left-handed fermions with quantum numbers as in Table III. The SM color is identified as the diagonal SU(3) in the $SU(3)\times SU(3)'$. The pNGB Higgs boson is coming from $\langle QQ\rangle\neq 0$, that spontaneously breaks $SU(5)\times U(1)'$ to SO(5). The second condensate that may form is $\langle\tilde{\chi}\chi\rangle$ that, if nonvanishing on the vacuum, would break $SU(3)\times SU(3)'\to SU(3)_c$. In the broken phase, no global anomalies are present, as the unbroken SU(3) corresponds to the gauged color. However, in the vacuum $\langle\tilde{\chi}\chi\rangle=0$, we need to match $A_{(SU(3)^3)}$ and $A_{(SU(3)'^3)}$ anomalies. In the underlying theory, $A_{(SU(3)^3)}=-A_{(SU(3)'^3)}=4$. All three-fermion bound states of the low energy theory [46] are of the form $\chi Q\chi$, and they transform as fundamental or two-index antisymmetric of the SU(3) flavor symmetries, and thus a matching of the anomaly seems possible. However, as they all contain a single Q, they all come with a multiplicity of 5, which means that in the low energy theory we get an anomaly coefficient multiple of 5 and not 4. We leave as an open question the role of bound states with more than three fermions, as $QQQ\chi\chi$, which in principle can transform as $\mathbf{3}$ or $\bar{\mathbf{3}}$ of one of the two global SU(3). These conclusions can be generalized to models with an arbitrary number of fermions, as the presence of a pNGB Higgs in the spectrum always requires that $N_Q\geq 5$.

The second model we analyze was proposed in Ref. [47]: it is based on a $G_{HC}=SU(3)$ gauge theory with seven vectorlike fermions. The SM gauge interactions are embedded in the diagonal SU(7) flavor symmetry, while top partners arise as three-fermion bound states, like baryons in QCD. A pNGB Higgs can be obtained in a similar way as in the more minimal coset $SU(4)^2\to SU(4)$ [11]. Setting aside the SM gauge interactions, it is a theory with a $SU(7)\times SU(7)\times U(1)$ global symmetry, broken in the vacuum by a condensate to the diagonal SU(7), and this

TABLE III. Fermionic field content of the theory of Ref. [46] and transformation under the global symmetry $SU(5)\times SU(3)\times SU(3)'\times U(1)^2$.

	SU(4)	SU(5)	$SU(3)\times SU(3)'$	$U(1)_X$	U(1)'
Q	**6**	**5**	$(\mathbf{1},\mathbf{1})$	0	−1
χ	**4**	**1**	$(\mathbf{3},\mathbf{1})$	−1/3	5/3
$\tilde{\chi}$	$\bar{\mathbf{4}}$	**1**	$(\mathbf{1},\bar{\mathbf{3}})$	1/3	5/3

does not have anomalies to be matched since it is a vectorlike theory. We believe unlikely that the strong dynamics breaks only partially the flavor symmetry to the diagonal $\mathrm{SU}(n)$ leaving a chiral $\mathrm{SU}(7-n)\times\mathrm{SU}(7-n)$ unbroken, because strong effects are flavor blind. Thus, in this theory no global anomaly can enforce composite fermions to be massless.

IV. CONCLUSIONS

In the context of composite pNGB Higgs models with top partial compositeness, we studied conditions under which the low energy theory might contain light top partners, lighter than the estimate based on naive dimensional analysis. The 't Hooft anomaly matching condition is proposed as a mechanism to force some fermions to be massless, or parametrically light once an explicit soft breaking of the global symmetry is introduced. We find this idea attractive because we are not aware of any other mechanism with the same effect, in four dimensions without elementary scalars, and because light top partners play an important role in model building and in new physics phenomenology at the LHC.

We examined in details specific models based on an UV free gauge theory of interacting fermions, proposed in the literature as underlying theories of pNGB Higgs with top partners, where we find no obstruction for this mechanism; namely the model possesses a global symmetry of which the anomaly can be matched by composite fermions with the quantum numbers of a top partner candidate. For this mechanism to work, we need at least two species of fundamental fermions since for one of them a global symmetry should survive while the other has to provide a condensate. This mechanism, however, does not apply to any UV construction; out of the three models under scrutiny, we found that solutions of the 't Hooft anomalies are only possible in one scenario, provided that one of the two fermion condensates does not occur. In the other two cases, either there are no solutions to the anomaly matching or the anomalies vanish. We believe this mechanism can thus be a useful criterion to select interesting models to be further studied, although we leave a thorough classification to a future work. The main question that remains open is about the vacuum that the theory chose to live in, and finding an answer requires a study of the model on the lattice, or by means of other nonperturbative techniques.

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