UV completions of partial compositeness: the case for a SU(4) gauge group

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ABSTRACT: We present a model of partial compositeness arising as the IR limit of a SU(4) gauge theory with only fermionic matter. This group is one of the most promising ones among a handful of possible choices allowing a symmetry breaking pattern incorporating custodial symmetry and a top partner candidate, while retaining asymptotic freedom. It is favored for not giving rise to lepto-quarks or Landau poles in the SM gauge couplings. The minimal UV theory consists of five hyperfermions in the anti-symmetric representation and three in the fundamental and anti-fundamental. The IR theory is centered around the coset SU(5)/SO(5), with top partners in the fundamental of SO(5), giving rise to one composite fermion of electric charge 5/3, three of charge 2/3 and one of charge $-1/3$. Electro-Weak symmetry breaking occurs via top-quark-driven vacuum misalignment. The top quark mass is generated via the mechanism of partial compositeness, while the remaining fermions acquire a mass via a standard quadratic coupling to the Higgs. We compute the top and bottom quark mass matrix and the Electro-Weak currents of the composite fermions. The model does not give rise to unacceptably large deviations from the SM $Z \to b\bar{b}$ decay width.

KEYWORDS: Technicolor and Composite Models, Beyond Standard Model

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Contents

1 Introduction

The discovery [1, 2] of a 126 GeV Higgs boson [3–5], together with our expectations from effective field theory, points to the existence of new states and enlarged symmetries at the LHC scale. While nowadays some degree of fine tuning seems almost unavoidable in any incarnation of this idea, due to the fierce direct and indirect experimental constraints, one possibility that still remains is the existence of a new strongly coupled gauge theory at a scale much below the GUT scale.

For this idea still to be viable today, some specific dynamical mechanisms must occur. Among the few possibilities, we concentrate on the following scenario, generally known under the name of "partial compositeness":

  i) First, the Higgs boson arises as a pseudo Nambu-Goldstone boson (pNGB) of a broken global symmetry and condenses at the EW scale $v = 246$ GeV via a "misalignment" mechanism [6]. This guarantees that the corrections to the S parameter are suppressed by a factor $v^2/f^2 \ll 1$, with f the decay constant of the pNGB.

  ii) Second, the top quark, (and possibly other fermions), acquires a mass by mixing with a composite state of the same quantum numbers [7]. This helps in suppressing flavor changing neutral currents (FCNC) and CP violating terms without reintroducing a large fine-tuning.

Many works on this subject start with a phenomenological lagrangian with the desired properties and use the CCWZ formalism [8, 9] to describe the interactions. Attempts to derive this lagrangian from an underlying model have been mostly based on the idea of extra dimensions and holography. We will not discuss these approaches in this paper and instead will refer to the many reviews [10–19] and references therein for the original literature. (We have been mostly following [20].)

Work on purely four-dimensional UV completions, based on some strongly coupled "hypercolor" (HC) group, has been hampered by the objective difficulty of constructing entirely satisfactory models giving rise to the two dynamical mechanisms above. One difficulty is in obtaining viable partners to all the Standard Model (SM) fermions. Another difficulty is in achieving realistic masses for those that do have a partner. One must require a mixing, schematically of the type $\bar{q}\mathcal{O}$, between a generic SM fermion q and a composite state $\mathcal{O}$. In order for this mechanism to be effective, the scaling dimension of $\mathcal{O}$ must be close to 5/2. This is easy to realize in the presence of elementary scalars ϕ in the HC theory as $\mathcal{O} \approx \phi\psi$ (where ψ is a HC fermion), but the reappearance of scalars calls once again for an explanation. This strategy is being pursued in the context of supersymmetric theories in e.g. [21, 22].

Purely fermionic UV completions require $\mathcal{O}$ to attain a large anomalous dimension. Apart for the exceptional case of an adjoint HC fermion ψ that can combine with the HC field strength F to give $\mathcal{O} \approx F_{\mu\nu}\gamma^{\mu\nu}\psi$ of perturbative dimension 7/2, the other possibility, for generic irreps, is to have some HC invariant combination $\mathcal{O} \approx \psi_1\psi_2\psi_3$ of perturbative dimension 9/2, requiring an anomalous dimension $\eta \approx -2$. While this is a tall order, it is nevertheless more appealing than the corresponding requirement needed for the pNGB composite operator $\mathcal{H}$ in the case where SM fermion masses are obtained by a bilinear term $\bar{q}\mathcal{H}q$. In this latter case [23], (see also [24–27][1]), the requirement on $\mathcal{H}$ is that it has scaling dimension close to 1, but this is the free field limit for a boson and implies that the scaling dimension of $\mathcal{H}^\dagger\mathcal{H}$ cannot be much different from 2, reintroducing the fine-tuning problem [28, 29]. On the contrary, 5/2 is safely above the free field case for a fermion and in any case it does not give rise to additional relevant perturbations. However, the idea [23] may still be viable for the SM fermions other than the top quark and we will rely on this in our construction.

A purely fermionic model of this type was proposed in [30] based on a HC group Sp(4) and some of its basic dynamical properties were studied. In [31] we classified, purely on group theoretical grounds, the models that fulfill the requirements i) and ii) above, together with some extra simplifying conditions such as a simple HC group. In [31] we made no attempt to study the dynamics of these models. In this work, we return to this issue and consider one of the most attractive models in the classification [31], based on a HC group SU(4).

Given that, in the most favorable possible scenario, the LHC will find evidence for compositeness that can be fully described by the IR effective theory, what is the interest

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[1]In particular, in [26, 27], some higher dimensional irreps have been studied that will also appear in the present work.

in looking for UV completions now? One reason is that, in the strictly IR approach, one has no control over the possible group realizations, (both the coset and the irreps) of the theory and one is forced to guess or to scan over "group theory space" (see e.g. [32, 33]). The UV completion can help pointing towards the most promising models. Equivalently, by considering what generic properties arise in the IR from a class of UV theories, one can test or rule out the whole UV class.

Let us summarize the organization and the main results of the paper.

In section 2, we present the UV theory. We discuss its matter content, the pattern of symmetry breaking and the composition of the top partners in terms of the hyperfermions. We show that the theory does not give rise to leptoquarks or any scalar composite state in the triplet or sextet of color. We compute the modification to the SM β-functions and show that no Landau pole arises at low scales. There is an amusing coincidence where the SM couplings almost unify but the scale at which it occurs is too small to be taken seriously and, at any rate, we know that new physics must arise before that to generate the needed couplings between the SM and the hypercolor sector.

Section 3 discusses the IR theory. We present the pNGB and top partner field content and argue that EW breaking proceeds as required. We then construct the relevant couplings between the SM and the composite fields. Due to the lack of potential partners for all SM fermions, partial compositeness is applied only to the top quark, and we propose that the remaining fields should be given a mass by standard quadratic interactions. We discuss what spurions should be used for this purpose. We construct the EW currents and the derivative couplings of the composite fermions. Here we find a happy circumstance when it comes to the $Z \to b\bar{b}$ decay. The irreps involved are such that the decay is safe from large corrections [34] arising from the composite partners. We also show this explicitly by going to the b mass eigenstates.

In section 4, we conclude with a short discussion and briefly review the current experimental status.

The main omission in this work is that we do not attempt to show that the anomalous dimensions for the composite operators are sufficient to realize a realistic mass spectrum, although arguments in favor of this possibility have been recently proposed in [30] for a similar model. Convincing evidence on this issue can only come via lattice simulations or a detailed analysis of the OPE that is beyond the scope of this paper. We also do not speculate on what physics could give rise to the required four-fermi couplings at a much higher scale.

2   The UV theory

In [31] we searched for gauge theories with fermionic matter allowing a spontaneous global symmetry breaking pattern $G_{\mathrm{F}}/H_{\mathrm{F}}$ compatible with custodial symmetry: $H_{\mathrm{F}} \supset G_{\mathrm{cus.}} \supset G_{\mathrm{SM}}$. (Having defined $G_{\mathrm{cus.}} = \mathrm{SU(3)}_c \times \mathrm{SU(2)}_L \times \mathrm{SU(2)}_R \times \mathrm{U(1)}_X$ and $G_{\mathrm{SM}} = \mathrm{SU(3)}_c \times \mathrm{SU(2)}_L \times \mathrm{U(1)}_Y$.) We further required the presence of one Higgs doublet $G_{\mathrm{F}}/H_{\mathrm{F}} \ni (\mathbf{1}, \mathbf{2}, \mathbf{2})_0$ of $G_{\mathrm{cus.}}$ and a composite fermionic trilinear partner for at least the third generation G_{SM} fermions $Q_L \in (\mathbf{3}, \mathbf{2})_{1/6}$ and $t_R \in (\mathbf{3}, \mathbf{1})_{2/3}$.

	SU(4)	SU(5)	SU(3)	SU(3)′	U(1)$_X$	U(1)′
ψ	**6**	**5**	**1**	**1**	0	−1
χ	**4**	**1**	**3**	**1**	−1/3	5/3
$\tilde\chi$	**$\bar4$**	**1**	**1**	**$\bar3$**	1/3	5/3

The top of the table has braces labelling the column groups: G_{HC} over SU(4), and G_{F} over the remaining columns.

Table 1. The fermions of the UV theory studied in this paper. They are to be thought of as two-component left-handed objects. Later, when discussing the low energy phenomenological lagrangian, we will find it more convenient to revert to four-component notation. G_{HC} is the hypercolor gauge group and G_{F} the global symmetry group before symmetry breaking.

We restricted the search to asymptotically free theories with a simple HC group G_{HC} and at most three inequivalent types of fermionic irreps. One could enlarge the class of theories, but the restricted class above already captures all the desired features. The solutions to the constraints above where presented in tables 2 and 3 of [31] and included the model presented in [30]. One can classify these models in various way. One possibility is to divide them according to the breaking of the global symmetry giving rise to the pNGB's. The two custodial cosets arising contain either SU(n)/Sp(n) or SU(n)/SO(n), with $n = 4$ and 5 being the minimal choice respectively.

Another distinction that can be made between them is whether they allow for composite scalars in the **3** or **6** of SU(3)$_c$. According to ones expectations, these are either exciting predictions or potential problems for these models and their role has been discussed in e.g. [35]. The model [30] contains such states originating from fermionic bilinears in the theory. We chose to work with theories that *do not* give rise to such states and this restricts the number of solutions considerably. In fact, this requirement, together with the requirement that the new hyperfermions do not give rise to Landau poles too close to the EW scale, essentially singles out one solution, presented here in table 1, based on the hypercolor group $G_{\mathrm{HC}} = $ SU(4) which will be the focus of this paper. It is interesting to notice that SU(4) is the only unitary group allowing this construction.

2.1 Field content of the UV theory

Since the full matter content in table 1 is non-chiral (**6** is a real irrep of SU(4)), the theory is manifestly free of gauge anomalies G_{HC}^3. The group $G_{\mathrm{F}} = $ SU(5) × SU(3) × SU(3)′ × U(1)$_X$ × U(1)′ describes the flavor group free of ABJ anomalies $G_{\mathrm{F}}G_{\mathrm{HC}}^2$. The QCD color gauge group SU(3)$_c$ is realized as the diagonal subgroup of SU(3) × SU(3)′, in perfect analogy with the flavor symmetries for the light quarks. The subgroup of G_{F} that does not give rise to 't Hooft anomalies G_{F}^3, and thus can be weakly gauged when coupled to the SM, is $H_{\mathrm{F}} = $ SO(5) × SU(3)$_c$ × U(1)$_X$, containing the custodial group $G_{\mathrm{cus.}}$ defined above.

The reason why it is not possible to build mesons (of any spin) bilinear in the $\chi, \tilde\chi$ fields and transforming in the **3** or **6** of SU(3)$_c$ is that χ and $\tilde\chi$ transform under a complex irrep of G_{HC}. In the present case, with the field content of table 1, after reducing the (**3**, **$\bar3$**)

Object	$SU(2)_L \times SU(2)_R$	$SU(5) \times SU(3) \times SU(3)' \times U(1)_X \times U(1)'$
$\chi\psi\chi$	$(1/2, 0)$	$(\mathbf{5}, \bar{\mathbf{3}}, \mathbf{1})_{(-2/3,\ 7/3)}$
	$(1/2, 0)$	$(\mathbf{5}, \mathbf{6}, \mathbf{1})_{(-2/3,\ 7/3)}$
	$(3/2, 0)$	$(\mathbf{5}, \mathbf{6}, \mathbf{1})_{(-2/3,\ 7/3)}$
$\tilde{\chi}\psi\tilde{\chi}$	$(1/2, 0)$	$(\mathbf{5}, \mathbf{1}, \mathbf{3})_{(2/3,\ 7/3)}$
	$(1/2, 0)$	$(\mathbf{5}, \mathbf{1}, \bar{\mathbf{6}})_{(2/3,\ 7/3)}$
	$(3/2, 0)$	$(\mathbf{5}, \mathbf{1}, \bar{\mathbf{6}})_{(2/3,\ 7/3)}$
$\bar{\tilde{\chi}}\bar{\psi}\chi$	$(1/2, 0)$	$(\bar{\mathbf{5}}, \mathbf{3}, \mathbf{3})_{(-2/3,\ 1)}$
	$(1/2, 1)$	$(\bar{\mathbf{5}}, \mathbf{3}, \mathbf{3})_{(-2/3,\ 1)}$
$\bar{\chi}\bar{\psi}\tilde{\chi}$	$(1/2, 0)$	$(\bar{\mathbf{5}}, \bar{\mathbf{3}}, \bar{\mathbf{3}})_{(2/3,\ 1)}$
	$(1/2, 1)$	$(\bar{\mathbf{5}}, \bar{\mathbf{3}}, \bar{\mathbf{3}})_{(2/3,\ 1)}$
$\bar{\chi}\psi\bar{\chi}$	$(1/2, 0)$	$(\mathbf{5}, \mathbf{3}, \mathbf{1})_{(2/3,\ -13/3)}$
	$(1/2, 1)$	$(\mathbf{5}, \bar{\mathbf{6}}, \mathbf{1})_{(2/3,\ -13/3)}$
$\bar{\tilde{\chi}}\psi\bar{\tilde{\chi}}$	$(1/2, 0)$	$(\mathbf{5}, \mathbf{1}, \bar{\mathbf{3}})_{(-2/3,\ -13/3)}$
	$(1/2, 1)$	$(\mathbf{5}, \mathbf{1}, \mathbf{6})_{(-2/3,\ -13/3)}$

Table 2. The composite fermionic operators classified according to their Lorentz and flavor quantum numbers. For each operator there is a corresponding conjugate one. After symmetry breaking, they combine into vector-like operators that create spin $1/2$ or spin $3/2$ resonances out of the vacuum.

of $SU(3) \times SU(3)'$ to color $SU(3)_c$, one can only construct color singlet/octet scalars[2] of type $\tilde{\chi}\chi$, $\tilde{\chi}^\dagger\chi^\dagger$, or color singlet/octet vectors $\chi^\dagger\chi$, $\tilde{\chi}^\dagger\tilde{\chi}$.

In table 2 we list all fermionic G_{HC} invariant that can be constructed with three elementary fields, together with their spin and G_F flavor quantum numbers (later to be broken to H_F). This list includes the top quark partners that will be of interest in the remaining sections.

2.2 Symmetry breaking in the UV theory

Now we would like to argue that the pattern of symmetry breaking to be expected in this model is $G_F \to H_F$, with G_F and H_F given above, while leaving the hypercolor gauge group $G_{HC} = SU(4)$ unbroken. Since ψ is in a real representation of G_{HC}, all the fermionic objects in table 2 can be made massive by giving a mass to the ψ fields. This means that none would be available to cancel the 't Hooft anomalies [36] associated to the G_F/H_F generators, which should then be broken [37]. This patter of symmetry breaking is also consistent with the arguments of [38].

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[2]We will sometimes drop all indices to avoid cluttering the formulas when the contractions are obvious. For instance, the vector octet is the traceless part of $\chi^\dagger_{ma}\sigma^\mu\chi^{mb}$. In the paper, $m, n \ldots, I, J \ldots, a, b \ldots$ and $a', b' \ldots$ are $SU(4)$, $SU(5)$, $SU(3)$ and $SU(3)'$ indices respectively and the contraction over the Weyl indices is understood.

A more dynamical argument is an adaptation of the Nambu-Jona-Lasinio method [39, 40] as done in [30]. The G_{HC}-invariant scalar bilinears that can be constructed are $\epsilon_{mnpq}\psi^{Imn}\psi^{Jpq}$, $\tilde{\chi}_{ma'}\chi^{ma}$ and their complex conjugates.

Introducing two auxiliary fields $M^{IJ} \equiv M^{JI}$ and $N_{a'}^a$, the fourth-order effective lagrangian can be written as

$$\mathcal{L} \supset -\frac{3}{2k}M^{IJ}M_{IJ}^\dagger - \frac{1}{2}M^{IJ}\epsilon^{mnpq}\psi_{Imn}^\dagger\psi_{Jpq}^\dagger - \frac{1}{2}M_{IJ}^\dagger\epsilon_{mnpq}\psi^{Imn}\psi^{Jpq} \qquad (2.1)$$
$$-\frac{9}{k'}N_a^{\dagger a'}N_{a'}^a - N_{a'}^a\tilde{\chi}^{\dagger ma'}\chi_{ma}^\dagger - N_a^{\dagger a'}\tilde{\chi}_{ma'}\chi^{ma}$$

which, eliminating M and N, reduces to

$$\mathcal{L} \supset \frac{k}{6}\epsilon_{mnpq}\psi^{Imn}\psi^{Jpq}\epsilon^{m'n'p'q'}\psi_{Im'n'}^\dagger\psi_{Jp'q'}^\dagger + \frac{k'}{9}\tilde{\chi}_{ma'}\chi^{ma}\tilde{\chi}^{\dagger na'}\chi_{na}^\dagger. \qquad (2.2)$$

The fields M and N can be reduced to non-negative diagonal form by Takagi and singular-value decomposition respectively

$$M^{IJ} = \sum_K \mu_K \Omega_K^I \Omega_K^J, \qquad\qquad \psi^K = \Omega_I^K \psi'^I \qquad\qquad (2.3)$$
$$N_{a'}^a = \sum_b \nu_b \tilde{\Xi}_{a'}^b \Xi_b^a, \qquad\qquad \chi^a = \Xi_b^a \chi'^b \qquad \text{and} \qquad \tilde{\chi}_{a'} = \tilde{\Xi}_{a'}^b \tilde{\chi}_b.$$

In (2.3), Ω, Ξ and $\tilde{\Xi}$ are orthogonal matrices and the sum is indicated explicitly only when the contraction is not manifestly group-covariant. Using (2.3), (2.1) becomes

$$\mathcal{L} \supset \sum_{I=1}^5 -\frac{3}{2k}\mu_I^2 - \frac{1}{2}\mu_I\epsilon^{mnpq}\psi_{Imn}'^\dagger\psi_{Ipq}'^\dagger - \frac{1}{2}\mu_I\epsilon_{mnpq}\psi'^{Imn}\psi'^{Ipq} \qquad (2.4)$$
$$\sum_{a=1}^3 -\frac{9}{k'}\nu_a^2 - \nu_a\tilde{\chi}'^{\dagger ma}\chi_{ma}'^\dagger - \nu_a\tilde{\chi}_{ma}'\chi'^{ma}.$$

Integrating out the fermions, with Λ interpreted as the G_{HC} strong scale, gives

$$\mathcal{L} \supset -\sum_{I=1}^5 V(\mu_I) - \sum_{a=1}^3 \mathrm{U}(\nu_a) \qquad (2.5)$$

with (using the same sharp cut-off as in [30] for simplicity)

$$V(\mu) = \frac{3}{2k}\mu^2 - \frac{3}{8\pi^2}\left(\Lambda^2\mu^2 + \Lambda^4\log\frac{\Lambda^2+\mu^2}{\Lambda^2} + \mu^4\log\frac{\mu^2}{\Lambda^2+\mu^2}\right)$$
$$\mathrm{U}(\nu) = \frac{9}{k'}\nu^2 - \frac{1}{2\pi^2}\left(\Lambda^2\nu^2 + \Lambda^4\log\frac{\Lambda^2+\nu^2}{\Lambda^2} + \nu^4\log\frac{\nu^2}{\Lambda^2+\nu^2}\right) \qquad (2.6)$$

A plot of the potential $V(\mu)$ is shown in figure 1. (The potential $\mathrm{U}(\nu)$ is obviously similar.) For large enough values of k the minimum is attained at non-zero μ and the symmetry is broken. This is not a proof of symmetry breaking since we have no control on the actual values of the effective couplings. It does however point to the same direction
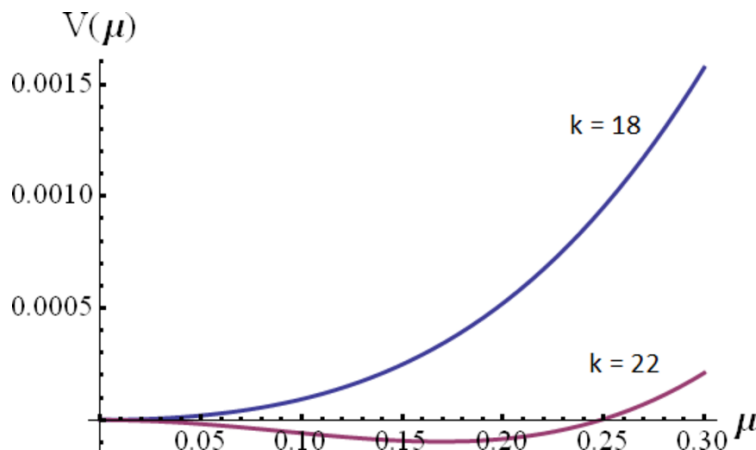
Figure 1. Plot of the potential $V(\mu)$ in units $\Lambda = 1$ for $k = 18$ and $k = 22$. The transition occurs at $k_{\min} \approx 20$. For $k > k_{\max} \approx 64$, the minimum is at a value comparable to the cutoff and the approximation breaks down.

as the previous argument and shows explicitly that, if symmetry breaking occurs, there is a basis in which $\langle \epsilon_{mnpq} \psi^{Imn} \psi^{Jpq} \rangle \propto \delta^{IJ}$, breaking $\mathrm{SU}(5) \to \mathrm{SO}(5)$, and $\tilde{\chi}_{ma'} \chi^{ma} \propto \delta^a_{a'}$, breaking $\mathrm{SU}(3) \times \mathrm{SU}(3)' \to \mathrm{SU}(3)_c$. ($\mathrm{U}(1)'$ is also broken while $\mathrm{U}(1)_X$ is left unbroken.)

The Maximally Attractive Channel hypotheses (MAC) [41] indicates that the breaking of $\mathrm{SU}(5)$ occurs at a higher scale compared to that of $\mathrm{SU}(3) \times \mathrm{SU}(3)'$. We can quantify the ratio of scales by a naive one-loop matching.

For the condensation in the $\psi\psi$ channel, the MAC is $\mathbf{6} \times \mathbf{6} \to \mathbf{1}$, with attractive strength $r_{\psi\psi} = C(\mathbf{1}) - C(\mathbf{6}) - C(\mathbf{6}) = -5$. ($C(\mathbf{R})$ is the quadratic Casimir of the irrep $\mathbf{R}$.) In the $\tilde{\chi}\chi$ channel we have a MAC $\bar{\mathbf{4}} \times \mathbf{4} \to \mathbf{1}$ and strength $C(\mathbf{1}) - C(\bar{\mathbf{4}}) - C(\mathbf{4}) = -15/4$. (The chiral channels like $\psi\chi$ are always less attractive than both of the above.)

The one loop $\mathrm{SU}(4)$ β-function coefficient with the ψ removed is $b = -38/3$, having defined $(\mu \mathrm{d}/\mathrm{d}\mu)\alpha_{\mathrm{HC}} = b\,\alpha^2_{\mathrm{HC}}/2\pi$. Assuming that the condensates form when $|r|\alpha_{\mathrm{HC}} \approx 1$, we can relate the scales as

$$\frac{\Lambda_{\psi\psi}}{\Lambda_{\tilde{\chi}\chi}} \approx \exp\left\{ \frac{2\pi}{|b|}(|r_{\psi\psi}| - |r_{\tilde{\chi}\chi}|) \right\} \approx 1.9 \tag{2.7}$$

Again, none of these arguments is rigorous (see e.g. [42, 43]) but it seems safe to assume that the $\mathrm{SU}(5)$ breaking occurs at a higher scale. We shall be mostly concentrating on the $\mathrm{SU}(5)/\mathrm{SO}(5)$ part, since this is where the EW dynamics resides. The effect of the remaining $\mathrm{SU}(3) \times \mathrm{SU}(3)'/\mathrm{SU}(3)_c$ is just that of generating a color octet pNGB that couples in the obvious way. We denote by f and f' the respective decay constants.

2.3 Running of the SM couplings

Having a candidate UV theory at one disposal can also be used to analyze the impact of the extra matter fields on the unification of the SM coupling. We should not expect any exact matching, since we have introduced a new gauge group and the new fields do not form complete multiplets. Morover, there is clearly some UV physics at higher scales
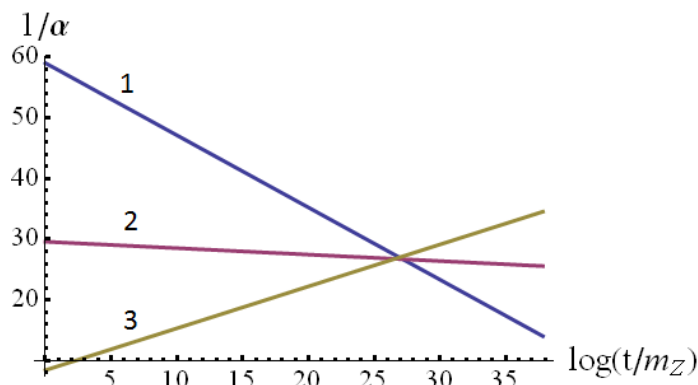
Figure 2. Running of the couplings in the present model. The matching is amusing but should not be taken seriously as it stands since it corresponds to a scale of $4.4 \ 10^{13}$ GeV.

still missing in order to explain the origin of the couplings between the hyperfermions and the SM fermions. At least though, one should check that the picture is not completely distorted, e.g. by the existence of Landau poles at low energies. In figure 2 we present the one-loop running of the SM couplings $\alpha_3 \equiv \alpha_s$, $\alpha_2 \equiv \alpha_W$ and $\alpha_1 \equiv 5\alpha_Y/3$ for our model. The running is given by the equation

$$\frac{\mathrm{d}}{\mathrm{d}t}\alpha_i^{-1} = -\frac{b_i}{2\pi}, \quad \text{with} \ \ t = \log(\mu/m_Z) \tag{2.8}$$

with[3]

$$(b_1, b_2, b_3) = (112/15, 2/3, -13/3). \tag{2.9}$$

It is amusing to see that the matching is improved, but this should not be taken seriously since the scale of the matching is way to small ($4.4 \ 10^{13}$ GeV) for proton decay to be ignored. Perhaps the addition of the extra heavy fields that are necessary anyway to generate the four-fermi interactions could cure that. We checked some of the remaining models classified in [31] and did not find any sign of unification. In fact, many suffer of problems from Landau poles.

3   The IR theory

Having discussed the basic features of the UV theory, we now present the IR effective theory. The two sets of fields that we will retain are the pNGBs and the top quark partners.

As before, we denote by Λ the SU(4) strong coupling scale, or, equivalently, the mass scale of a *typical* composite state (i.e. neither Goldstone not the lightest top-like partner, that will be assumed to have lower mass M). f is the SU(5)/SO(5) pNGB decay constant. The ratio between Λ and f can be estimated as $\Lambda/f \equiv g \approx 4\pi/\sqrt{N_{\rm HC}} = 2\pi$.

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[3]For comparison, we recall the well known results $(b_1, b_2, b_3) = (41/10, -19/6, -7)$ and $(b_1, b_2, b_3) = (33/5, 1, -3)$ for the SM and MSSM respectively as well as $(b_1, b_2, b_3) = (152/15, -2, -11/3)$ for the model [30]. The Reduced Planck scale corresponds to $t = 37.8$.

For guidance, a tuning parameter $\xi \equiv v^2/f^2 \approx 0.1$ gives $f \approx 800\,\mathrm{GeV}$ and $\Lambda \approx 5\,\mathrm{TeV}$. The top-partner mass M will lie somewhere in between f and Λ. The UV description in terms of the SU(4) gauge theory coupled to the SM is assumed to be valid up to a UV scale $\Lambda_{\mathrm{UV}} \gg \Lambda$ where the four-fermion interactions are generated. We will not attempt to discuss the physics involved at Λ_{UV}, but FCNC indicate that $\Lambda_{\mathrm{UV}} > 10^7\,\mathrm{GeV}$.

3.1 Composite fields

As far as the pNGBs are concerned, we argued in the previous section that the symmetry breaking pattern takes the form

$$\frac{G_{\mathrm{F}}}{H_{\mathrm{F}}} = \frac{\mathrm{SU}(5) \times \mathrm{SU}(3) \times \mathrm{SU}(3)' \times \mathrm{U}(1)_X \times \mathrm{U}(1)'}{\mathrm{SO}(5) \times \mathrm{SU}(3)_c \times \mathrm{U}(1)_X}$$

$$= \left(\frac{\mathrm{SU}(5)}{\mathrm{SO}(5)}\right) \times \left(\frac{\mathrm{SU}(3) \times \mathrm{SU}(3)'}{\mathrm{SU}(3)_c}\right) \times \mathrm{U}(1)' \tag{3.1}$$

So far, all these bosons are massless and we now need to discuss how the coupling to the SM fields affects their spectrum.

The Goldstone boson η' corresponding to $\mathrm{U}(1)'$ is SM-neutral and will remain massless in our approximation. We will drop it from now on, but its role should be discussed in the cosmological context.

The EW breaking will be driven by the dynamics of the SU(5)/SO(5) coset and for this we need to specify the embedding of the EW part of the SM gauge group $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ into H_{F}. This is done by first decomposing $\mathrm{SO}(5) \to \mathrm{SO}(4) \equiv \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R$, then identifying a $\mathrm{U}(1)_R$ subgroup of $\mathrm{SU}(2)_R$ generated by T_R^3 and, lastly, setting $Y = T_R^3 + X$.

We take the vev for the ψ bilinear $\langle \psi^I \psi^J \rangle$ proportional to δ^{IJ}.[4] The 24 traceless hermitian generators of the fundamental irrep of SU(5) are then decomposed into 10 imaginary anti-symmetric generators of SO(5), generically denoted by T^a, and the remaining 14 traceless real symmetric broken generators, generically denoted by S^i, corresponding each to one Goldstone boson. The generators of SO(4) are embedded into SO(5) by padding the last row and column with zeros and choosing the remaining 4×4 representation as in [32].[5] It is convenient to have the expression for the generators of the $\mathrm{SU}(2)_L \times \mathrm{U}(1)_R$ subgroup of SO(5):

$$T_L^1 = \frac{i}{2}\begin{pmatrix} 0 & 0 & 0 & -1 & 0 \\ 0 & 0 & -1 & 0 & 0 \\ 0 & 1 & 0 & 0 & 0 \\ 1 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 \end{pmatrix}, \qquad T_L^2 = \frac{i}{2}\begin{pmatrix} 0 & 0 & 1 & 0 & 0 \\ 0 & 0 & 0 & -1 & 0 \\ -1 & 0 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 \end{pmatrix}$$

$$T_L^3 = \frac{i}{2}\begin{pmatrix} 0 & -1 & 0 & 0 & 0 \\ 1 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & -1 & 0 \\ 0 & 0 & 1 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 \end{pmatrix}, \qquad T_R^3 = \frac{i}{2}\begin{pmatrix} 0 & -1 & 0 & 0 & 0 \\ 1 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 1 & 0 \\ 0 & 0 & -1 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 \end{pmatrix} \tag{3.2}$$

---

[4]This is different from the most commonly used convention established in [44–46], where this coset was originally presented. The physical couplings are of course independent on the explicit representation chosen.

[5]Much of the notation in this work is influenced by this paper.

As far as the broken generators are concerned, we describe them by giving the explicit matrix for the Goldstone fields. Decomposing the SU(5)/SO(5) coset according to SU(2)$_L \times$ U(1)$_R$ one finds [44] one totally SM neutral real boson η (that will also be dropped in the following), one "true Higgs" doublet H, a Y-neutral, SU(2)$_L$-triplet Φ_0 and a charged one $\Phi_\pm$:

$$\mathbf{14} \to \mathbf{1}_0 + \mathbf{2}_{\pm 1/2} + \mathbf{3}_0 + \mathbf{3}_{\pm 1} \equiv (\eta, H, \Phi_0, \Phi_\pm) \tag{3.3}$$

For the Higgs, we will use the standard notation $H = (H_+, H_0)$, the indices denoting directly the electric charge $Q = T_L^3 + Y \equiv T_L^3 + T_R^3$ (recall that all Goldstone bosons have $X = 0$). For the triplets we use the notation $\phi_{m_R}^{m_L}$, with $m_{R/L} = -1, 0, +1$ eigenvalues of $T_{R/L}^3$, e.g. $\Phi_0 \supset (\phi_0^-, \phi_0^0, \phi_0^+)$, and $\Phi_+ \supset (\phi_+^-, \phi_+^0, \phi_+^+)$. The electric charge is $Q = m_L + m_R$ and $(\phi_{m_R}^{m_L})^\dagger = \phi_{-m_R}^{-m_L}$. There is thus [46] one double-charge scalar (ϕ_+^+ and h.c.), two single-charge ones (ϕ_+^0, ϕ_0^+ and h.c.) and three neutral ones, (having dropped η and η'), (ϕ_0^0, $\Re\phi_+^-$ and $\Im\phi_+^-$). All these pNGB fit into the SU(5)/SO(5) generators as

$$H = \begin{pmatrix} 0 & 0 & 0 & 0 & -iH_+/\sqrt{2} \\ 0 & 0 & 0 & 0 & H_+/\sqrt{2} \\ 0 & 0 & 0 & 0 & iH_0/\sqrt{2} \\ 0 & 0 & 0 & 0 & H_0/\sqrt{2} \\ -iH_+/\sqrt{2} & H_+/\sqrt{2} & iH_0/\sqrt{2} & H_0/\sqrt{2} & 0 \end{pmatrix}$$

$$\Phi_0 = \begin{pmatrix} \phi_0^0/\sqrt{2} & 0 & i(\phi_0^- - \phi_0^+)/2 & (\phi_0^- + \phi_0^+)/2 & 0 \\ 0 & \phi_0^0/\sqrt{2} & (\phi_0^- + \phi_0^+)/2 & -i(\phi_0^- - \phi_0^+)/2 & 0 \\ i(\phi_0^- - \phi_0^+)/2 & (\phi_0^- + \phi_0^+)/2 & -\phi_0^0/\sqrt{2} & 0 & 0 \\ (\phi_0^- + \phi_0^+)/2 & -i(\phi_0^- - \phi_0^+)/2 & 0 & -\phi_0^0/\sqrt{2} & 0 \\ 0 & 0 & 0 & 0 & 0 \end{pmatrix}$$

$$\Phi_+ = \begin{pmatrix} \phi_+^+/\sqrt{2} & i\phi_+^+/\sqrt{2} & i\phi_+^0/2 & \phi_+^0/2 & 0 \\ i\phi_+^+/\sqrt{2} & -\phi_+^+/\sqrt{2} & -\phi_+^0/2 & i\phi_+^0/2 & 0 \\ i\phi_+^0/2 & -\phi_+^0/2 & \phi_+^-/\sqrt{2} & -i\phi_+^-/\sqrt{2} & 0 \\ \phi_+^0/2 & i\phi_+^0/2 & -i\phi_+^-/\sqrt{2} & -\phi_+^-/\sqrt{2} & 0 \\ 0 & 0 & 0 & 0 & 0 \end{pmatrix} \tag{3.4}$$

We combine these bosons as

$$\Pi = H + H^\dagger + \Phi_0 + \Phi_+ + \Phi_+^\dagger, \quad \text{and} \quad \Sigma = \exp\left(\frac{i\Pi}{f}\right) \tag{3.5}$$

Note that Π is a real and symmetric matrix. We will later argue that EW breaking takes place as expected, namely by giving a vev to the neutral component of H, $H_0 = h/\sqrt{2}$. The remaining components of H are then the true Goldstone bosons to be eaten by the $W^\pm$ and Z and can be set to zero in the unitary gauge.[6] It is convenient to express the

---

[6]Throughout the paper, we use the normalization where the W mass, the vev of the canonically normalized h, and the decay constant f are related by $m_{\rm w} = (g/2)f \sin(\langle h \rangle/f)$, i.e. *the same* relation as used in the smaller coset SO(5)/SO(4), yielding $v = f \sin(\langle h \rangle/f) = 246\,{\rm GeV}$. We find this uniform definition less confusing than the one more commonly used for this coset, where f is scaled by a factor 2.

Object			SO(5) × SU(3)$_c$ × U(1)$_X$
$\tilde{\chi}\psi\tilde{\chi}$,	$\bar{\chi}\psi\bar{\chi}$,	$2 \times \bar{\chi}\bar{\psi}\tilde{\chi}$	$(\mathbf{5}, \mathbf{3})_{2/3}$
$\chi\psi\chi$,	$\bar{\tilde{\chi}}\psi\bar{\tilde{\chi}}$,	$2 \times \bar{\tilde{\chi}}\bar{\psi}\chi$	$(\mathbf{5}, \bar{\mathbf{3}})_{-2/3}$

Table 3. The spin 1/2 color triplets particles created by composite fermions after symmetry breaking. Shown are the LH combinations. The charge conjugates of the operators in the second row combine with the ones in the first row to give a total of four Dirac spinors. We assume without proof that one of them is significantly lighter than the others. Similar considerations can be made for the sextets and the spin 3/2 resonances, although in this case we don't need to assume that some of them are lighter than the generic scale Λ.

exponential exactly in h and expand around the other fields, if necessary, using

$$\Sigma = \Sigma(h) + \frac{i}{f} \int_0^1 \mathrm{d}s \Sigma((1-s)h) \left(\Phi_0 + \Phi_+ + \Phi_+^\dagger \right) \Sigma(sh) + \dots \tag{3.6}$$

where, defining $c_h = \cos(h/f)$ and $s_h = \sin(h/f)$,

$$\Sigma(h) = \begin{pmatrix} 1 & 0 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 & 0 \\ 0 & 0 & 1 & 0 & 0 \\ 0 & 0 & 0 & c_h & is_h \\ 0 & 0 & 0 & is_h & c_h \end{pmatrix} \tag{3.7}$$

The Goldstone bosons in the SU(3) × SU(3)′/SU(3)$_c$ coset transform in the **8** of color. We simply write them as $\pi = \pi^a \lambda^a / 2$ where λ^a are the usual Gell-Mann matrices.

Moving on to the top quark partners, one of the advantages of having a candidate UV completion is that it allows one to motivate picking a particular irrep of H_F for such objects. We collect in table 3 all the spin-half SU(3)$_c$ triplet excitations created by the composite states obtained from table 2, now classified according to the unbroken global symmetry.

We identify both the **5** and $\bar{\mathbf{5}}$ of SU(5) with the **5** of SO(5) and construct the SU(3)$_c$ irreps from SU(3) × SU(3)′ using $\mathbf{3} \times \mathbf{3} = \bar{\mathbf{3}} + \mathbf{6}$. We do not consider any longer the spin 3/2 objects nor the color sextets, that we assume correspond to heavier states at the scale Λ. These states however are additional prediction of this UV completion and would allow one to discern it from other possibilities if experiments were performed at a higher scale.

The breaking of the global symmetry is what allows us to form Dirac fields out of the LH objects displayed in table 3 and their RH conjugates.

One assumption (that we will not attempt to justify in this work) is that one linear combination of operators creates a fermionic resonance of mass M that is lighter than the generic resonance scale Λ. This is not too unreasonable since we are asking for less that a factor ten suppression. Thus, we will assume the existence of *one Dirac field* Ψ, of mass M, transforming in the $(\mathbf{5}, \mathbf{3})_{2/3}$ of H_F.

To extract the EW quantum numbers for these fields, note that

$$
\begin{array}{ccc}
\mathrm{SO}(5) \times \mathrm{SU}(3)_c \times \mathrm{U}(1)_X & \qquad & (\mathbf{5},\mathbf{3})_{2/3} \\
\downarrow & & \downarrow \\
G_{\mathrm{cus.}} \equiv \mathrm{SU}(3)_c \times \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_X & & (\mathbf{3},\mathbf{2},\mathbf{2})_{2/3} + (\mathbf{3},\mathbf{1},\mathbf{1})_{2/3} \\
\downarrow & & \downarrow \\
G_{\mathrm{SM}} \equiv \mathrm{SU}(3)_c \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y & & (\mathbf{3},\mathbf{2})_{7/6} + (\mathbf{3},\mathbf{2})_{1/6} + (\mathbf{3},\mathbf{1})_{2/3} \\
\downarrow & & \downarrow \\
\mathrm{SU}(3)_c \times \mathrm{U}(1)_{\mathrm{e.m.}} & & \mathbf{3}_{5/3} + 3 \times \mathbf{3}_{2/3} + \mathbf{3}_{-1/3}
\end{array}
\tag{3.8}
$$

We have thus succeeded in contracting a partner for the LH SM field $q_L^3 = (t_L, b_L)$, namely the RH projection of the Dirac field $(T, B) \in (\mathbf{3},\mathbf{2})_{1/6}$ above, and a partner to the RH SM field t_R, namely the LH projection of the Dirac field $R \in (\mathbf{3},\mathbf{1})_{2/3}$. We will only consider the mixing between the composite fermions and the third family.

Notice that we do not find partners to the remaining SM fields, including b_R. We could simply ignore this problem by focusing on the more pressing issue of the top mass, but we argue that for the remaining particles it is still feasible to consider a bilinear mass term as in the early constructions [6] and we will do so in the following. Given the quantitative difference of the top quark mass and the difficulty in finding an acceptable model giving all partners, this option seems more attractive to us.

We have already discussed the components T, B and R. Denoting the remaining fields by $(X, Y) \in (\mathbf{3},\mathbf{2})_{7/6}$, we write the full $(\mathbf{5},\mathbf{3})_{2/3}$ multiplet as

$$
\Psi = \frac{1}{\sqrt{2}} \begin{pmatrix} iB - iX \\ B + X \\ iT + iY \\ -T + Y \\ \sqrt{2}iR \end{pmatrix}
\tag{3.9}
$$

After EW symmetry breaking, the fields T, Y and R acquire electric charge $2/3$ and mix with the top quark. Similarly, the field B acquires an electric charge $-1/3$ and mixes with the bottom quark. The field X has charge $5/3$ and is a generic prediction of many models of this type.

3.2 EW symmetry breaking

The most pressing issue is to show that the desired misalignment of the Higgs field H, leading to the correct EW symmetry breaking, can occur.

Precise quantitative computations are precluded by our lack of control of the strong dynamics. What we can hope to do is to show that the couplings of the SM fields to the pNGBs are such that the misalignment can occur for the Higgs doublet H but not for the other fields. We will consider top-quark-driven misalignment as proposed in [47].

We want to write an effective action coupling the pNGBs to the SM vector bosons and fermions. Under a generic global $g \in \mathrm{SU}(5)$ transformation, Σ in (3.5) transforms non-linearly as $\Sigma \to g\Sigma h^T(\Pi, g)$ with $h(\Pi, g) \in \mathrm{SO}(5)$, a real matrix. In this case, we

are allowed to construct a simpler object $U = \Sigma\Sigma^T \equiv \exp\left(\frac{2i}{f}\Pi\right)$ transforming linearly as $U \to gUg^T \equiv \mathrm{Symm}_g \circ U$.

The couplings to the vector bosons are obtained from

$$\mathcal{L} \supset \frac{f^2}{16}\mathrm{tr}\left((D_\mu U)^\dagger D^\mu U\right) \tag{3.10}$$

where,

$$D_\mu U = \partial_\mu U - igW_\mu^a[T_L^a, U] - ig'B_\mu[T_R^3, U]. \tag{3.11}$$

For simplicity, we will only consider contributions from the $SU(2)_L$ EW bosons W_μ^a.

To couple the pNGBs to the SM fermions we need to determine the spurionic embeddings[7] by considering how they can be coupled to the composite field Ψ. Given Ψ in the **5** of SO(5) as above, we can construct the operators $\Sigma\Psi$ and $\Sigma^*\Psi$ transforming in the **5** and **$\bar{5}$** of SU(5) respectively. This fact forces us to pick, as spurionic embedding of the elementary quarks q_L and t_R, the **5** and **$\bar{5}$** representation as well. We write

$$\hat{q}_L = \frac{1}{\sqrt{2}}\begin{pmatrix} ib_L \\ b_L \\ it_L \\ -t_L \\ 0 \end{pmatrix}, \quad \text{and} \quad \hat{t}_R = \begin{pmatrix} 0 \\ 0 \\ 0 \\ 0 \\ it_R \end{pmatrix} \tag{3.12}$$

The coupling with the Ψ will be important later, now we focus on the induced terms. In momentum space they read

$$\mathcal{L} \supset G(p)\left(\bar{\hat{q}}_L U\hat{t}_R + \bar{\hat{t}}_R U^*\hat{q}_L\right) \tag{3.13}$$

where $G(p)$ is a form factor depending on the strong dynamics and the rest of the fields is evaluated at zero momentum. Notice that kinetic terms of type $\bar{\hat{q}}_L U \not{p}\hat{q}_L$ are *not* allowed since U is in the **15**. For the same reason, we cannot pick both the spurions in the **5**.

We start by expanding around the *unbroken* vacuum $\Pi = 0$ and look for possible destabilizing effects. Once we convince ourselves that the breaking occurs when H_0 gets a vev, we set all other fields to zero and treat H_0 to all orders.

The couplings in (3.10) with the $SU(2)_L$ EW bosons is proportional to

$$\mathrm{tr}(T_L^a T_L^a \Pi\Pi - T_L^a \Pi T_L^a \Pi) = \frac{3}{2}H^\dagger H + 4\Phi_+^\dagger\Phi_+ + 2\Phi_0^\dagger\Phi_0 \tag{3.14}$$

Vector couplings of this type do not misalign the vacuum [48]. This means that they will contribute to the pNGB potential with a positive overall coefficient to the combination in (3.14). The only possible negative contributions must come from the fermionic couplings, which are proportional to

$$\bar{\hat{q}}_L\Pi\hat{t}_R + \bar{\hat{t}}_R\Pi\hat{q}_L = \frac{2}{f}(\bar{q}_L H^\dagger t_R - \bar{t}_R H q_L). \tag{3.15}$$

---

[7]For a given SM field q, we denote by $\hat{q}$ a field with the same dynamical content but transforming as a full multiplet of G_F. For conciseness, we call the whole $\hat{q}$ "spurion" and never write down the auxiliary fields.

(a) Contribution of the SU(2)$_L$ gauge bosons.

(b) Contribution of the top quark.
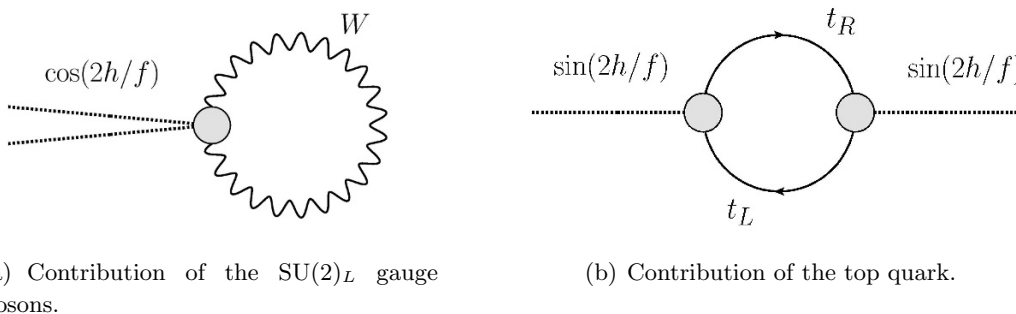
Figure 3. The leading order contributions to the induced Higgs potential.

Hence, it is only for the field H that we can expect a misalignment. We now set $H_0 = h/\sqrt{2}$, all other fields to zero, and write

$$\mathrm{U}(h) = \Sigma(h)\Sigma(h)^T \equiv \Sigma(2h) \tag{3.16}$$

yielding

$$W_\mu^a W_\mu^b \mathrm{tr}(\mathrm{U}(h)T_L^a \mathrm{U}(h)^\dagger T_L^b) = \frac{1}{2}\left(1 + \cos(2h/f)\right) W_\mu^c W_\mu^c$$

$$\bar{\hat{q}}_L \mathrm{U}(h)\hat{t}_R + \bar{\hat{t}}_R \mathrm{U}(h)^* \hat{q}_L = \frac{1}{\sqrt{2}}\sin(2h/f)(\bar{t}_L t_R + \bar{t}_R t_L). \tag{3.17}$$

The contribution to the Coleman-Weinberg potential [49] is given, to leading order, by the diagrams in figure 3. We can then summarize the contribution of the integral over the resonances of the strong sector by two dimensionless numbers α and β as done in e.g. [20]

$$V(h) \propto \alpha \cos(2h/f) - \beta \sin^2(2h/f). \tag{3.18}$$

An acceptable EW breaking minimum will be attained for $\beta \gtrsim |\alpha|/2$ at $\sin^2(2\langle h\rangle/f) = 1 - (\alpha/2\beta)^2$. Recalling that with our conventions $v = f\sin(\langle h\rangle/f)$, we get a relation between the fine-tuning parameter ξ and the terms in the Higgs potential

$$\xi \equiv \left(\frac{v}{f}\right)^2 \approx \frac{1}{4}\left(1 - \left(\frac{\alpha}{2\beta}\right)^2\right), \tag{3.19}$$

i.e. a factor of four improvement over the minimal coset. We believe it makes sense to compare the two because the relation between v, $\langle h\rangle$ and f has been chosen to be the same for both. To our knowledge, this last observation was first made in [50], but see e.g. [51–53] for related recent work.

A simpler analysis can be done to show that the color octet π^a will not be destabilized and thus color remains unbroken. These pNGBs will gain a mass that we can roughly estimate from the diagram in figure 4 as

$$m_\pi^2 \approx \frac{3 \times 8 \times \alpha_s}{4\pi}\Lambda^2 \approx (2\ \mathrm{TeV})^2 \tag{3.20}$$
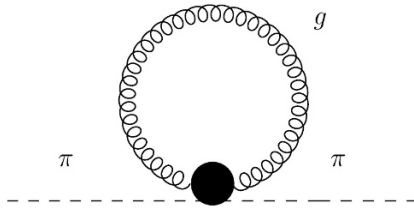
Figure 4. The gluon contribution to the mass of the pNGB π.

	$SU(3)_c$	$U(1)_X$	$SU(2)_L$	$U(1)_R$
$\hat{q}_L^i$	**3**	$2/3$	**2**	$-1/2$
$\hat{u}_R^i$	**3**	$2/3$	**1**	0
$\hat{d}_R^i$	**3**	$2/3$	**1**	-1

Table 4. The remaining spurions quantum numbers.

In (3.20), 3 is the number of off-shell gluon polarizations, 8 the color factor and Λ^2 summarizes the effects of the strong sector resonances. One could be more precise and use sum-rules to write this contribution in terms of the masses of the lowest lying states but not much is gained in this case since we do not have experimental information on their mass. In this case there can also be contributions from the quarks kinetic terms but we expect them to be subleading compared to the gluons.

3.3 The fermionic mass terms

In subsection 3.1 we discussed the transformation properties of the composite pNGB and fermions. We saw that, in order to couple linearly to the top quark, we needed to embed the elementary fields q_L and t_R into spurions in the **5** and $\bar{\mathbf{5}}$ of SU(5) respectively. However, the lack of candidate partners for the remaining fermions made it impossible to construct all masses this way. The complexity (and unlikeliness) of UV completion providing all partners made us propose a compromise: the remaining masses are constructed with *bilinears*. At least the fine tuning is mitigated by only needing to achieve a mass of the order of a few GeV. Even this requires some care though, since, in order to preserve the SU(5) invariance, we need to pick the representation for the spurions in a way compatible with the choices in subsection 3.1.

Let us consider the quarks and denote by $\hat{q}_L^i$, $\hat{u}_R^i$, and $\hat{d}_R^i$, the SU(5) spurionic embeddings, where $i = 1, 2, 3$ is the family number. We have already encountered $\hat{q}_L$ (no index i) and $\hat{t}_R$ which have the same physical field content as $\hat{q}_L^3$ and $\hat{u}_R^3$ but arranged in a different irrep.

Since the pNGB fields carry no $U(1)_X$ charge and the quarks mix, the remaining spurions must have the same $U(1)_X$ charge as the top partners. This fact, together with the usual relation between X, Y and T_R^3, fixes the quantum numbers displayed in table 4. Now we want to embed the $SU(2)_L \times U(1)_R$ irreps of table 4 into irreps of SU(5) that allow to construct mass terms $\bar{\hat{q}}_L^i U^* \hat{u}_R^j$ and $\bar{\hat{q}}_L^i U \hat{d}_R^j$. One sees that, restricting to at most

"two-index" irreps, the only solution that allows reproducing table 4 and constructing the needed mass terms is $\hat{q}_L^i \in \mathbf{24}$, $\hat{u}_R^i \in \mathbf{10}$, and $\hat{\tilde{d}}_R^i \in \bar{\mathbf{10}}$, the adjoint, anti-symmetric and conjugate anti-symmetric irreps respectively.[8] Setting all the auxiliary fields to zero we have, explicitly

$$
\hat{q}_L^i = \frac{1}{2}
\begin{pmatrix}
0 & 0 & 0 & 0 & id_L^i \\
0 & 0 & 0 & 0 & d_L^i \\
0 & 0 & 0 & 0 & iu_L^i \\
0 & 0 & 0 & 0 & -u_L^i \\
id_L^i & d_L^i & iu_L^i & -u_L^i & 0
\end{pmatrix},
\qquad
\hat{u}_R^i = \frac{1}{2}
\begin{pmatrix}
0 & u_R^i & 0 & 0 & 0 \\
-u_R^i & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & -u_R^i & 0 \\
0 & 0 & u_R^i & 0 & 0 \\
0 & 0 & 0 & 0 & 0
\end{pmatrix},
$$

$$
\hat{\tilde{d}}_R^i = \frac{1}{2\sqrt{2}}
\begin{pmatrix}
0 & 0 & id_R^i & -d_R^i & 0 \\
0 & 0 & d_R^i & id_R^i & 0 \\
-id_R^i & -d_R^i & 0 & 0 & 0 \\
d_R^i & -id_R^i & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0
\end{pmatrix}.
\tag{3.21}
$$

The same construction works for the leptons with the only difference that now the $\mathrm{U}(1)_X$ charges are taken to be zero.

We are now ready to write down the mass terms for the top and bottom quarks, including the contribution to $b_R \in \hat{\tilde{d}}_R^3$. We concentrate on the third family and write

$$
\mathcal{L} \supset \frac{M}{2}\bar{\Psi}\Psi + \lambda_q f \bar{\hat{q}}_L \Sigma \Psi_R + \lambda_t f \bar{\hat{t}}_R \Sigma^* \Psi_L + \sqrt{2}\mu_b \, \mathrm{tr}\left(\bar{\hat{q}}_L^3 U \hat{\tilde{d}}_R^3\right) + \mathrm{h.c.}
\tag{3.22}
$$

The dimensionless couplings λ_q and λ_t between SM fields and the composite fermion are expected to be of the same order and control the mass of the top quark. The mass parameter μ_b is required to give a mass to the bottom quark and we ignore the subleading quadratic terms like $\bar{\hat{q}}_L U \hat{t}_R$ and $\mathrm{tr}\left(\bar{\hat{q}}_L^3 U^* \hat{u}_R^3\right)$ for the top.

Inserting (3.16) and (3.21) in (3.22), we obtain the following mass matrices

$$
\mathcal{M}_T =
\begin{pmatrix}
0 & \frac{\lambda_q}{2}f(1+\mathrm{c}_h) & \frac{\lambda_q}{2}f(1-\mathrm{c}_h) & \frac{\lambda_q}{\sqrt{2}}f\mathrm{s}_h \\
\frac{\lambda_t}{\sqrt{2}}f\mathrm{s}_h & M & 0 & 0 \\
-\frac{\lambda_t}{\sqrt{2}}f\mathrm{s}_h & 0 & M & 0 \\
\lambda_t f\mathrm{c}_h & 0 & 0 & M
\end{pmatrix}
\tag{3.23}
$$

and

$$
\mathcal{M}_B =
\begin{pmatrix}
\mu_b \mathrm{s}_h \mathrm{c}_h & \lambda_q f \\
0 & M
\end{pmatrix}
\tag{3.24}
$$

in terms of which the lagrangian (3.22) can be written as

$$
\mathcal{L} \supset (\bar{t}_L, \bar{T}_L, \bar{Y}_L, \bar{R}_L) \cdot \mathcal{M}_T \cdot
\begin{pmatrix}
t_R \\
T_R \\
Y_R \\
R_R
\end{pmatrix}
+ (\bar{b}_L, \bar{B}_L) \cdot \mathcal{M}_B \cdot
\begin{pmatrix}
b_R \\
B_R
\end{pmatrix}
+ \mathrm{h.c.}
\tag{3.25}
$$

---

[8]The $\mathbf{10}$ and $\bar{\mathbf{10}}$ irreps can be used interchangeably. That the fundamental irrep cannot be used can be inferred by the need to reproduce $T_R^3 = -1$ for $\hat{\tilde{d}}_R^3$. The spurions are normalized to have canonical kinetic energy, e.g. $\mathrm{tr}(\bar{\hat{q}}_L^i \not{\partial} \hat{q}_L^i) = \bar{q}_L^i \not{\partial} q_L^i$.

The lowest singular values of the two mass matrices have to be made coincide with the known masses of the top and bottom quarks. For the top sector, we expand the lowest mass to leading order in the higgs vev v to find

$$m_t \approx \frac{\sqrt{2} M f \lambda_q \lambda_t}{\sqrt{M^2 + \lambda_q^2 f^2} \sqrt{M^2 + \lambda_t^2 f^2}} \, v, \tag{3.26}$$

proportional to the product $\lambda_q \lambda_t$, in agreement with diagrammatic expectations. The remaining masses have non-vanishing values even for $v \to 0$ and are, to zeroth order, equal to M, $\sqrt{M^2 + \lambda_q^2 f^2}$ and $\sqrt{M^2 + \lambda_t^2 f^2}$

For the bottom quark we can go to the mass eigenstates by writing

$$\begin{pmatrix} b_L \\ B_L \end{pmatrix} = \begin{pmatrix} \cos \lambda & \sin \lambda \\ -\sin \lambda & \cos \lambda \end{pmatrix} \begin{pmatrix} b'_L \\ B'_L \end{pmatrix} \equiv R_\lambda \begin{pmatrix} b'_L \\ B'_L \end{pmatrix}$$

$$\begin{pmatrix} b_R \\ B_R \end{pmatrix} = \begin{pmatrix} \cos \rho & \sin \rho \\ -\sin \rho & \cos \rho \end{pmatrix} \begin{pmatrix} b'_R \\ B'_R \end{pmatrix} \equiv R_\rho \begin{pmatrix} b'_R \\ B'_R \end{pmatrix} \tag{3.27}$$

with

$$\tan 2\rho = \frac{2\mu_b f \lambda_q s_h c_h}{M^2 + \lambda_q^2 f^2 - \mu_b^2 s_h^2 c_h^2}, \quad \text{and} \quad \tan 2\lambda = \frac{2\lambda_q f M}{M^2 - \lambda_q^2 f^2 - \mu_b^2 s_h^2 c_h^2}. \tag{3.28}$$

The mass of the b quark is, to lowest order in the Higgs vev,

$$m_b \approx \frac{\mu_b M}{f \sqrt{M^2 + \lambda_q^2 f^2}} \, v \tag{3.29}$$

and the remaining mass is equal, to zeroth order, to $\sqrt{M^2 + \lambda_q^2 f^2}$, thus nearly degenerate with one of the top partners.

3.4 The fermionic currents

We now compute the contribution of the fermionic partners to the EW currents. For this, we need first to define the matrix-valued one-forms

$$p_\mu = \sum_{i=1}^{14} S^i \text{tr}\big(S^i \Sigma^{-1}(i\partial_\mu \Sigma + e V_\mu \Sigma) \big), \quad v_\mu = \sum_{a=1}^{10} T^a \text{tr}\big(T^a \Sigma^{-1}(i\partial_\mu \Sigma + e V_\mu \Sigma) \big) \tag{3.30}$$

where

$$V_\mu = W_\mu^+ \frac{T_L^1 + i T_L^2}{\sqrt{2} s_w} + W_\mu^- \frac{T_L^1 - i T_L^2}{\sqrt{2} s_w} + \left(A_\mu + \frac{c_w}{s_w} Z_\mu \right) T_L^3 + \left(A_\mu - \frac{s_w}{c_w} Z_\mu \right) T_R^3 \tag{3.31}$$

is the matrix-valued SM gauge field and e is the electric coupling constant.

The gauge currents associated to the composite fermion can be read off from

$$\mathcal{L} \supset \bar{\Psi} \gamma^\mu \left(\frac{2}{3} e A_\mu \mathbf{1} - \frac{2 s_w}{3 c_w} e Z_\mu \mathbf{1} + v_\mu \right) \Psi + K \bar{\Psi} \gamma^\mu p_\mu \Psi$$

$$\equiv e(J_A^\mu A_\mu + J_Z^\mu Z_\mu + J_{W^+}^\mu W_\mu^+ + J_{W^-}^\mu W_\mu^-) + \dots \tag{3.32}$$

where the only undetermined constant is K and the dots represent terms without gauge fields. Extracting the coefficients, we obtain, restricting to the coupling with the pNGB h only,

$$
\begin{aligned}
J_Z^\mu &= C_{XX}\bar{X}\gamma^\mu X + C_{TT}\bar{T}\gamma^\mu T + C_{YY}\bar{Y}\gamma^\mu Y + C_{RR}\bar{R}\gamma^\mu R \\
&\quad + C_{BB}\bar{B}\gamma^\mu B + C_{RT}(\bar{R}\gamma^\mu T + \text{h.c.}) + C_{RY}(\bar{R}\gamma^\mu Y + \text{h.c.}) + C_{TY}(\bar{T}\gamma^\mu Y + \text{h.c.}) \\
J_{W^+}^\mu &= C_{XT}\bar{X}\gamma^\mu T + C_{XY}\bar{X}\gamma^\mu Y + C_{XR}\bar{X}\gamma^\mu R \\
&\quad + C_{TB}\bar{T}\gamma^\mu B + C_{YB}\bar{Y}\gamma^\mu B + C_{RB}\bar{R}\gamma^\mu B
\end{aligned}
\tag{3.33}
$$

and, of course, $J_{W^-}^\mu = (J_{W^+}^\mu)^\dagger$ and

$$
J_A^\mu = \frac{5}{3}\bar{X}\gamma^\mu X + \frac{2}{3}(\bar{T}\gamma^\mu T + \bar{Y}\gamma^\mu Y + \bar{R}\gamma^\mu R) - \frac{1}{3}\bar{B}\gamma^\mu B.
\tag{3.34}
$$

The coefficients in (3.33) are given by

$$
\begin{aligned}
C_{XX} &= \frac{1}{\text{s}_\text{w}\text{c}_\text{w}}\left(\frac{1}{2} - \frac{5\text{s}_\text{w}^2}{3}\right) \\
C_{TT} &= -\frac{2\text{s}_\text{w}}{3\text{c}_\text{w}} + \frac{\text{c}_h}{2\text{s}_\text{w}\text{c}_\text{w}} \\
C_{YY} &= -\frac{2\text{s}_\text{w}}{3\text{c}_\text{w}} - \frac{\text{c}_h}{2\text{s}_\text{w}\text{c}_\text{w}} \\
C_{RR} &= -\frac{2\text{s}_\text{w}}{3\text{c}_\text{w}} \\
C_{BB} &= \frac{1}{\text{s}_\text{w}\text{c}_\text{w}}\left(-\frac{1}{2} + \frac{\text{s}_\text{w}^2}{3}\right) \\
C_{TY} &= 0 \\
C_{RT} = C_{RY} &= \frac{K\,\text{s}_h}{2\sqrt{2}\text{s}_\text{w}\text{c}_\text{w}} \\
C_{XT} = C_{YB} &= \frac{1 - \text{c}_h}{2\sqrt{2}\text{s}_\text{w}} \\
C_{XY} = C_{TB} &= \frac{1 + \text{c}_h}{2\sqrt{2}\text{s}_\text{w}} \\
C_{RB} = -C_{XR} &= \frac{K\,\text{s}_h}{2\text{s}_\text{w}}
\end{aligned}
\tag{3.35}
$$

The important point to notice is the value of eC_{BB}, which coincides with the analogous coefficient arising from the elementary b_L. This guarantees that, after rotating to the mass eigenbasis with the matrices (3.27), the branching ratio $Z \to b\bar{b}$ does not suffer large corrections. This is an explicit realization of the mechanism described in [34]. Here the situation is satisfactory since the UV completion has *forced us* to choose a b_L belonging to one of the "custodial irreps" described in [34].

3.5 Additional couplings

There are infinite series of additional couplings dictated by the non-linear structure of the Lagrangian. For instance, so far we have not considered the colored pNGBs, but their

interactions can be written down in analogy with three-flavor QCD, with the difference that now the "baryons" Ψ are in a triplet

$$\mathcal{L} \supset \frac{1}{2} \bar{\Psi} \gamma^\mu \left(g_s G_\mu^a + \frac{\lambda_A}{f'} \gamma^5 \partial_\mu \pi^a + \frac{1}{f'^2} f^{abc} \pi^b \partial_\mu \pi^c + \dots \right) \lambda^a \Psi. \tag{3.36}$$

(Here G_μ^a is the gluon, g_s the QCD coupling constant and λ_A the analog of the axial coupling.)

The term in (3.32) containing p_μ also gives rise to derivative interactions with the Higgs field of type

$$\mathcal{L} \supset \frac{iK}{\sqrt{2}f} (\bar{R} \gamma^\mu Y - \bar{R} \gamma^\mu T) \partial_\mu h + \text{h.c.} \tag{3.37}$$

as well as couplings with the other pNGBs. Among these, there are non-derivative $1/f$-suppressed couplings between two composite fermions, a gauge field and a pNGB that could also be relevant for phenomenology.

Lastly, the mixing between composite and elementary quarks in (3.22) also gives rise to couplings with the additional pNGBs. Here we present only those that survive the limit $\langle h \rangle / f \to 0$

$$\mathcal{L} \supset \lambda_q \left(\bar{b}_L Y_R \phi_-^0 - \bar{t}_L X_R \phi_-^0 - i\sqrt{2} \bar{b}_L X_R \phi_-^- + i\sqrt{2} \bar{t}_L Y_R \phi_-^+ + \frac{i}{\sqrt{2}} \bar{b}_L B_R \phi_0^0 \right.$$
$$\left. - \frac{i}{\sqrt{2}} \bar{t}_L T_R \phi_0^0 - \bar{b}_L T_R \phi_0^- + \bar{t}_L B_R \phi_0^+ \right) + \text{h.c.} \tag{3.38}$$

with no term arising in this limit from the couplings to the t_R.

4 Discussion

We presented a model of partial compositeness motivated by an UV completion based on a SU(4) gauge group. This group was shown in [31] to be the only unitary group allowing for custodial symmetry and top partners while retaining asymptotic freedom. It is also the preferred one if one wants to avoid scalar color triplets and sextets.

The fields in the IR theory carrying SM charges consists of the standard SU(5)/SO(5) pNGB coset, a color octet pNGB, one fermion mixing with the bottom quark, three mixing with the top quark and one of electric charge 5/3. The top quark gained mass purely via the mechanism of partial compositeness while, for the lighter fermions, we resorted to quadratic couplings because of the lack of potential partners.

Much remains to be done before this model can be considered fully satisfactory. The main issue is whether the dynamics of the gauge theory is such that a realistic mass spectrum can be justified. Here we are forced to work at small N_{HC}, so analysis similar to those in [24, 25] could turn out to be useful. Still, we felt that the IR theory is sufficiently appealing to motivate our study, and thus we presented the most important couplings and discussed the mass spectrum for the top and bottom sector.

The S and T constraints [54–56] for this class of models have been discussed in many places and reviewed in [20]. These contributions can be made acceptable at the cost of

tuning the parameter ξ to be sufficiently small. As for the $Z \to b\bar{b}$ decay, we showed that the model belongs to the class of models for which this channel is protected from acquiring large deviations from the SM result. (Top quark compositeness may also constrain these models, see [57] for an extensive discussion.)

The LHC direct searches during run 1 have probed a large fraction of these models. Limits on the fermionic partners have been set to $m_X \gtrsim 800\,\text{GeV}$, $m_T \gtrsim 700\,\text{GeV}$ and $m_B \gtrsim 700\,\text{GeV}$ for the charge 5/3 [58], charge 2/3 [59] and charge $-1/3$ [60] respectivelly.[9] Searches for doubly charged Higgs bosons appeared in [61] with the data at 7 TeV, setting a bound of $m_\phi \gtrsim 400\,\text{GeV}$ in various dilepton channels. The search for a generic scalar color octet (called s8 in [62, 63]) has excluded a mass range in the region between 1. and 2.66 TeV, but this limit needs to be analyzed carefully before applying it to the octet appearing in this paper. Lastly, the search for vector resonances [64, 65], also expected to appear at a scale Λ in the strongly coupled sector, has set a bound of $m_\rho \gtrsim 1.1\,\text{TeV}$. The next LHC run will probe even deeper into these classes of models, exceeding the TeV limit for all composite fermions [66] and ensuring plenty of entertainment for the coming years.

Acknowledgments

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