Large N scalars: From glueballs to dynamical Higgs models

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We construct effective Lagrangians, and corresponding counting schemes, valid to describe the dynamics of the lowest lying large N stable massive composite state emerging in strongly coupled theories. The large N counting rules can now be employed when computing quantum corrections via an effective Lagrangian description. The framework allows for systematic investigations of composite dynamics of a non-Goldstone nature. Relevant examples are the lightest glueball states emerging in any Yang-Mills theory. We further apply the effective approach and associated counting scheme to composite models at the electroweak scale. To illustrate the formalism we consider the possibility that the Higgs emerges as the lightest glueball of a new composite theory; the large N scalar meson in models of dynamical electroweak symmetry breaking; the large N pseudodilaton useful also for models of near-conformal dynamics. For each of these realizations we determine the leading N corrections to the electroweak precision parameters. The results nicely elucidate the underlying large N dynamics and can be used to confront first principle lattice results featuring composite scalars with a systematic effective approach.

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I. INTRODUCTION

Strong dynamics continues to pose a formidable challenge. Several ingenious analytical and numerical techniques have been invented, exploited, and are routinely used to elucidate some of its physical properties. The large number of underlying colors limit is a time-honored example [1–3]. It has been extensively used in quantum chromodynamics (QCD) and string theory, and it constitutes the backbone of the gauge-gravity duality program. We will further highlight here the power of the large N expansion by introducing a four-dimensional calculable framework permitting us to investigate the dynamics of the lightest stable non-Goldstone large N composite state.

’t Hooft and Witten demonstrated that Yang-Mills theories for a large number of colors admit an effective description in terms of an infinite number of noninteracting absolutely stable hadronic states of arbitrary spin [1–3]. By capitalizing on this central result we focus on the physics of the lightest massive scalar state that is known to play an important role in QCD [4–13] and in various models of dynamical electroweak symmetry breaking as summarized in [14]. The framework can also be used to consistently determine quantum corrections to compare with first principle lattice simulations of composite dynamics featuring scalars [15–18].

The paper is organized as follows. In Sec. II we introduce the effective theory for the lightest massive glueball scalar state emerging within a pure Yang-Mills theory, and provide the associated counting scheme. Here we discuss the intriguing interplay between momentum and large N expansions. The framework goes beyond the glueball example and lays the foundation of the subsequent analyses. We then extend the framework to several models of composite electroweak dynamics in Sec. III where the Higgs is identified with the lightest composite state. In Sec. IV we determine the large N dependence, for each model, of the electroweak precision parameters [19] stemming from different dynamical Higgs realizations. We summarize our results in Sec. V.

II. LARGE N EFFECTIVE THEORY FOR THE LIGHTEST GLUEBALL STATE

Consider the lightest scalar state stemming from an $SU(N)$ pure Yang-Mills theory, which is also expected to be the lowest lying state of the full theory. At an infinite number of colors the effective Lagrangian for this state is simply the one with a free scalar field [1–3],

$$L_{GB} = \frac{1}{2}\partial_\mu h \partial^\mu h - \frac{m_h^2}{2}h^2 + \mathcal{O}(N^{-1}). \tag{1}$$

It is possible to go beyond the free-field limit by first defining with Λ_H the intrinsic composite scale of the theory that permits us to expand the effective Lagrangian in both $1/N$ and ∂^2/Λ_H^2 as follows:

$$L_{GB} = \frac{1}{2}\partial_\mu h \partial^\mu h - \frac{m_h^2}{2}\sum_{q=0,p=0}^{\infty} V_{q,p}\left(\frac{\partial^2}{\Lambda_H^2}\right)^q \left(\frac{1}{N}\frac{h}{\Lambda_H}\right)^p h^2. \tag{2}$$

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Here $V_{q,p}$ are dimensionless coefficients of order unity with the tree-level values $V_{0,0} = 1$ and $V_{1,0} = 0$ ensuring canonically normalized mass and kinetic terms. The expansion in $1/N$ takes care of the large N suppression of higher point correlators, while the higher derivative terms take into account integrating out heavier states. The notation for the higher order derivative terms schematically indicates all possible Lorentz invariant operators one can construct, acting over the h field, for a given exponent q. To leading order in the double expansion we have

$$L_{GB} = \frac{1}{2}\partial_\mu h \partial^\mu h - \frac{m_h^2}{2}\left[1 + \frac{V_{0,1}}{N}\frac{h}{\Lambda_H}\right]h^2. \quad (3)$$

This shows that the trilinear coupling of the scalar is naturally suppressed in this limit. Expanding a little further we have

$$L_{GB} = \frac{1}{2}\partial_\mu h \partial^\mu h - \frac{m_h^2}{2}\left[1 + \frac{V_{0,1}}{N}\frac{h}{\Lambda_H} + \frac{V_{1,1}}{N}\frac{\partial^2}{\Lambda_H^2}\frac{h}{\Lambda_H} + \frac{V_{0,2}}{N^2}\frac{h^2}{\Lambda_H^2} + V_{2,0}\left(\frac{\partial^2}{\Lambda_H^2}\right)^2\right]h^2. \quad (4)$$

In our schematic representation of higher derivatives terms we have it, for example, that $\partial^2 h^3$ represents the operator $h\partial_\mu h \partial^\mu h$. The other superficially different operator containing the same number of fields and derivatives, i.e. $h^2 \Box h$, is not independent because it is related to $h\partial_\mu h\partial^\mu h$ via an integration by parts.

We have ordered the terms in such a way that the $1/N$ order counts as ∂^2/Λ_H^2.[1] However, in the following we will work in the very low momentum regime and therefore we drop the derivative terms and obtain

$$L_{GB} = \frac{1}{2}\partial_\mu h \partial^\mu h - \frac{m_h^2}{2}h^2\left[1 + \frac{V_{0,1}}{N}\frac{h}{\Lambda_H} + \frac{V_{0,2}}{N^2}\frac{h^2}{\Lambda_H^2}\right]. \quad (5)$$

The effective theory features small self-couplings, even though it stems from a highly nonperturbative underlying gauge theory. The glueball mass receives $1/N^2$ corrections at the fundamental theory level. By computing the one loop corrections to the h two-point function one can check that the effective theory correctly reproduces the expected large N corrections. Of course, the effective Lagrangian is not renormalizable in the usual sense but due to the $1/N$ and ∂^2/Λ_H^2 ordering it is possible to organize and subtract the divergences order by order in this double expansion. Furthermore the coefficients of this effective theory can be determined, for a given underlying Yang-Mills gauge theory, via lattice simulations [20].

[1]This means that the momentum expansion shell is taken to be of the same order as the $1/N$ expansion.

III. LARGE N SCALARS FOR DYNAMICAL HIGGS MODELS

We now extend the framework presented above in order to introduce consistent effective descriptions of dynamical Higgs models. We are not concerned with fitting the latest experimental data but focus instead on elucidating the associated large N dynamics and effective theory structure.

We shall first introduce different examples and then, for each of these examples, we compute the electroweak precision observables in Sec. IV, more specifically the S and T parameters [19].

A. The dynamical Higgs as the lightest large N glueball

We start by considering the logical possibility that the dynamical Higgs state is the lightest glueball of a new fundamental composite theory. Besides pure Yang-Mills one can also consider theories with matter. Let us consider as a concrete example an $SU(N)$ gauge theory with two Dirac flavors transforming according to the two-index representation of the underlying gauge theory. Unlike the case in which the underlying fermions transform according to the fundamental representation, like QCD, in the two-index case at a large number of colors the axial anomaly is not suppressed [21]. This implies that the eta prime does not become massless in the large N limit and, besides the pions of the theory, the lightest scalar of the theory behaves as a glueball [22]. In fact in these theories mesons and glueballs at large N have the same large N properties. One could also imagine theories displaying large distance conformality in which one adds explicit sources of conformal breaking, such as fermion masses. That the lightest states in this case are glueballs has been shown to occur via lattice simulations [23] and via controllable perturbative examples [24,25]. However, while in the first explicit example (the one with fermions in higher-dimensional representations) the large N limit is legitimate, in the near-conformal cases one has also to consider the large number of the flavor limit, which is beyond the scope of this work. Of course, in the calculable case of [24,25] there is no need to write an effective field theory.

Within this scenario one can envision the newly discovered particle at the Large Hadron Collider (LHC) to be the lightest glueball state of a new Yang-Mills theory with a new N-independent string tension proportional to Λ_H. The scale is not automatically related, here, with the electroweak symmetry breaking scale $v \simeq 246$ GeV or $4\pi v$. Therefore dynamical spontaneous breaking of the electroweak symmetry is triggered by either another strongly coupled sector or, if within the same theory, via a distinct dynamically induced chiral symmetry scale.[2] Since the state h is a singlet with respect to the standard model (SM) symmetries we can write

[2]We remind the reader that in theories with an intact center group the confining and the chiral scale are well separated [26].

$$\mathcal{L}_{\text{Glueball-Higgs}} = \mathcal{L}_{\overline{\text{SM}}} + \left(1 + \frac{2r_\pi}{N\Lambda_H}h + \frac{s_\pi}{N^2\Lambda_H^2}h^2\right)\frac{v^2}{4}\text{Tr}D_\mu U^\dagger D^\mu U + \frac{1}{2}\partial_\mu h \partial^\mu h - \frac{m_h^2}{2}h^2\left[1 + \frac{V_{0,1}}{N}\frac{h}{\Lambda_H} + \frac{V_{0,2}}{N^2}\frac{h^2}{\Lambda_H^2}\right]$$
$$- m_t\left(1 + \frac{r_t}{N\Lambda_H}h\right)\left[\bar{q}_L U\left(\frac{1}{2} + T^3\right)q_R + \text{H.c.}\right] - m_b\left(1 + \frac{r_b}{N\Lambda_H}h\right)\left[\bar{q}_L U\left(\frac{1}{2} - T^3\right)q_R + \text{H.c.}\right] + \cdots$$
$$+ \mathcal{O}\left(\frac{1}{4\pi v}, \frac{\partial^2}{\Lambda_H^2}\right), \tag{6}$$

where $\mathcal{L}_{\overline{\text{SM}}}$ is the SM Lagrangian without Higgs and Yukawa sectors, the ellipses denote Yukawa interactions for SM fermions other than the top-bottom doublet $q \equiv (t, b)$, and $\mathcal{O}(1/\Lambda_H)$ includes higher-dimensional operators, which are suppressed by powers of $1/\Lambda_H$. In this Lagrangian U is the usual exponential map of the Goldstone bosons produced by the breaking of the electroweak symmetry, $U = \exp(i2\pi^a T^a/v)$, with covariant derivative $D_\mu U \equiv \partial_\mu U - igW_\mu^a T^a U + ig' U B_\mu T^3$, $2T^a$ are the Pauli matrices, with $a = 1, 2, 3$. The kinetic term and the potential of the SM Higgs have been replaced by the effective theory for the lightest glueball state. The tree-level SM is recovered for $r_\pi = r_t = r_b = N\frac{\Lambda_H}{v}$ and $s_\pi = N^2\frac{\Lambda_H^2}{v^2}$. Here we will keep these couplings at order unity. Also we have not speculated on the hidden sector providing the link between the new glueball theory and the SM sector, but required it to respect the large N counting for the insertion of an extra glueball-Higgs interaction. We have also ordered the higher derivatives on h such that they are subleading when compared to the $1/N$ operators retained here.

If we consider the infinite number of colors limit of the new Yang-Mills theory we arrive first at a perturbative self-interacting glueball state coupled to the SM also via perturbative couplings. We have, therefore, at our disposal an organization structure that allows us to go beyond the tree level. We shall investigate the dependence on the number of new colors N in the section dedicated to the electroweak parameters.

B. The large N physics of the dynamical Higgs

In time-honored models of dynamical electroweak symmetry breaking [27,28] the Higgs can be identified with a fermion-antifermion meson.[3] Depending on the new fermion representation with respect to the underlying gauge group one can have different large N countings [14] such as the Corrigan and Ramond one [21]. The counting is incorporated directly into the pion decay constant $F_\Pi^2 = d(R)\Lambda_{TC}^2$ with $d(R)$ being the dimension of the technicolor theory and Λ_{TC} an intrinsic scale not dependent on the number of colors [29,30].

Let us assume for definitiveness that we have an $SU(\bar{N})$ underlying theory featuring a doublet of techniquarks transformed according to the representation R of the composite theory and therefore we can set $v = F_\Pi(\bar{N})$. For a generic N it is sufficient to replace v with $v\sqrt{d/\bar{d}}$. For example for the fundamental representation $d = N$ and $\bar{d} = \bar{N}$. Unlike the previous glueball-Higgs example here also the pion sector is affected by the large N scaling since the self-interactions among the composite pions are also controlled by N. Choosing for definitiveness the underlying fermions to belong to the fundamental representation we have

$$\mathcal{L}_{TC}(N) = \mathcal{L}_{\overline{\text{SM}}} + \frac{N}{\bar{N}}\left(1 + \frac{2r_\pi}{v}\sqrt{\frac{\bar{N}}{N}}h + \frac{\bar{N}}{N}\frac{s_\pi}{v^2}h^2\right)\frac{v^2}{4}\text{Tr}D_\mu U^\dagger D^\mu U + \frac{1}{2}\partial_\mu h\partial^\mu h - \frac{m_h^2}{2}h^2\left[1 + \frac{\sqrt{\bar{N}}V_{0,1}}{\sqrt{N}}\frac{h}{v} + \frac{\bar{N}V_{0,2}}{N}\frac{h^2}{v^2}\right]$$
$$- \frac{y_t v}{\sqrt{2\bar{N}}}\sqrt{N}\left(1 + \frac{\sqrt{\bar{N}}r_t}{\sqrt{N}v}h\right)\left[\bar{q}_L U\left(\frac{1}{2} + T^3\right)q_R + \text{H.c.}\right] - \frac{y_b v}{\sqrt{2\bar{N}}}\sqrt{N}\left(1 + \frac{\sqrt{\bar{N}}r_b}{\sqrt{N}v}h\right)\left[\bar{q}_L U\left(\frac{1}{2} - T^3\right)q_R + \text{H.c.}\right]$$
$$+ \cdots + \mathcal{O}\left(\frac{1}{4\pi v}, \frac{\partial^2}{v^2}\right). \tag{7}$$

We have therefore

$$m_W^2 = g^2 v^2 \frac{N}{4\bar{N}}, \qquad m_q = y_q v\sqrt{\frac{N}{2\bar{N}}}, \tag{8}$$

[3] This does not always have to be the case, meaning that the lightest state can be, in principle, made by multifermion states.

with q a given quark, g the weak coupling, and y_q the Yukawa coupling. The physical electroweak scale is now identified with the following combination:

$$\sqrt{\frac{N}{\bar{N}}}v = 2\frac{m_W}{g}. \tag{9}$$

In the effective Lagrangian we arrive at

$$\mathcal{L}_{TC}(N) = \mathcal{L}_{\overline{\rm SM}} + \left(1 + g\frac{2r_\pi}{2m_W}h + g^2\frac{s_\pi}{4m_W^2}h^2\right)\frac{m_W^2}{g^2}{\rm Tr}D_\mu U^\dagger D^\mu U + \frac{1}{2}\partial_\mu h\partial^\mu h - \frac{m_h^2}{2}h^2\left[1 + gV_{0,1}\frac{h}{2m_W} + g^2V_{0,2}\frac{h^2}{4m_W^2}\right]$$
$$- m_t\left(1 + g\frac{r_t}{2m_W}h\right)\left[\bar{q}_L U\left(\frac{1}{2} + T^3\right)q_R + {\rm H.c.}\right] - m_b\left(1 + g\frac{r_b}{2m_W}h\right)\left[\bar{q}_L U\left(\frac{1}{2} - T^3\right)q_R + {\rm H.c.}\right] + \cdots. \tag{10}$$

The $1/\sqrt{N}$ cost of introducing an extra power of the composite field h is monitored nicely by the corresponding power in g.

C. The large N dynamical pseudodilaton

It is possible to imagine that the Higgs state of the SM is associated with the spontaneous breaking of a conformal symmetry. There are several possible realisations according to which the breaking of the conformal dynamics can be associated with either a nonperturbative sector [31,32] or a perturbative one [24,25].

The large N counting and the request to satisfy the conformal relations can both be ensured by imposing

$$N\Lambda_H = f, \qquad r_q = r_\pi = s_\pi = 1, \tag{11}$$

with f an N-dependent scale. The Lagrangian in (12) then becomes

$$\mathcal{L}_{\rm Dilaton}(N) = \mathcal{L}_{\overline{\rm SM}} + \left(1 + \frac{2h}{N\Lambda_H} + \frac{h^2}{N^2\Lambda_H^2}\right)\frac{v^2}{4}{\rm Tr}D_\mu U^\dagger D^\mu U + \frac{1}{2}\partial_\mu h\partial^\mu h - \frac{m_h^2}{2}h^2\left[1 + \frac{V_{0,1}}{N}\frac{h}{\Lambda_H} + \frac{V_{0,2}}{N^2}\frac{h^2}{\Lambda_H^2}\right]$$
$$- m_t\left(1 + \frac{h}{N\Lambda_H}\right)\left[\bar{q}_L U\left(\frac{1}{2} + T^3\right)q_R + {\rm H.c.}\right] - m_b\left(1 + \frac{h}{N\Lambda_H}\right)\left[\bar{q}_L U\left(\frac{1}{2} - T^3\right)q_R + {\rm H.c.}\right] + \cdots. \tag{12}$$

This framework can be immediately extended to any dilatoniclike interpretation of the state h, such as the one coming from a near conformal-like technicolor dynamics where f (Λ_H) is identified with v. Because we would like to investigate the explicit N dependence we hold fixed the N-independent scale Λ_H. Clearly this limit corresponds to the glueball-Higgs case with extra constraints for the h couplings.

We are now ready to investigate the first consequences of the large N counting.

IV. S AND T PARAMETERS

As a relevant application of the formalism introduced above we study two important correlators, i.e. the S and T parameters [33] for the three different types of dynamical Higgs models discussed above.

Defining by S the difference between S in the full theory $S_{\rm theory}$, and the value of S in the SM, i.e. $S_{\rm SM}$, we arrive at [34]

$$S = \left(1 - \frac{\kappa_1^2}{4}\right)\left[\frac{f(m_Z/m_h)}{6\pi} + \frac{1}{12\pi}\log\frac{(4\pi\Lambda_H)^2}{m_h^2} + \frac{5}{72\pi}\right] + 16\pi c(4\pi\Lambda_H) + \frac{1}{12\pi}\log\frac{m_h^2}{m_{h,\rm ref}^2}$$
$$+ \frac{f(m_Z/m_{h,\rm ref}) - f(m_Z/m_h)}{6\pi}, \tag{13}$$

where κ_1 is the coefficient of the linear term in h/v multiplying the operator ${\rm Tr}[D_\mu U D^\mu U^\dagger]$, which in the SM is equal to two, and

$$f(x) \equiv \frac{2x^2 + x^4 - 3x^6 + (9x^4 + x^6)\log x}{(1-x^2)^3}. \quad (14)$$

The c term in (13) is a needed counterterm whose finite part will be reabsorbed in the definition of the cutoff scale and therefore will be set to zero in the following.

For the T parameter one obtains [34]

$$T = -\frac{3}{16\pi\cos^2\theta_w}\left(1 - \frac{\kappa_1^2}{4}\right)\log\frac{(4\pi\Lambda_H)^2}{m_h^2} + T_{\rm SM}^h(m_h) - T_{\rm SM}^h(m_{h,\rm ref}), \quad (15)$$

where, as for the S parameter, we have absorbed the finite part of the counterterm in the actual value of Λ_H. $m_{h,\rm ref}$ is the reference value of the Higgs mass, θ_w is the Weinberg angle, and $T_{\rm SM}^h(m_h)$ is given in Eq. (22) of [34].

We can now determine the large N behavior of these relevant parameters coming from the various large N dynamical Higgs models introduced earlier.

A. S and T for the large N glueball and dilaton-Higgs models

In the glueball case we have

$$\kappa_1^{GB} = 2\frac{v r_\pi}{N\Lambda_H}, \quad (16)$$

and the SM limit is recovered when $N\Lambda_H = v r_\pi$. The precision parameters are then

$$S_{GB} = \left(1 - \frac{v^2 r_\pi^2}{N^2\Lambda_H^2}\right)\left[\frac{f(m_Z/m_h)}{6\pi} + \frac{1}{12\pi}\log\frac{(4\pi\Lambda_H)^2}{m_h^2} + \frac{5}{72\pi}\right] + \frac{1}{12\pi}\log\frac{m_h^2}{m_{h,\rm ref}^2} + \frac{f(m_Z/m_{h,\rm ref}) - f(m_Z/m_h)}{6\pi}, \quad (17)$$

and

$$T_{GB} = -\frac{3}{16\pi\cos^2\theta_w}\left(1 - \frac{v^2 r_\pi^2}{N^2\Lambda_H^2}\right)\log\frac{(4\pi\Lambda_H)^2}{m_h^2} + T_{\rm SM}^h(m_h) - T_{\rm SM}^h(m_{h,\rm ref}). \quad (18)$$

We have chosen $m_h = 125$ GeV and $m_{h,\rm ref} = 117$ GeV. The first observation is that if the glueball-Higgs scale Λ_H is larger than the electroweak scale v there is a positive contribution to the S parameter and an associated negative one for the T parameter. On the other hand, we observe a reduction (increase) of the S (T) parameter if Λ_H is smaller than v. This is an intriguing general result given the fact that the scale of compositeness is $4\pi\Lambda_H$ can be kept above the electroweak scale.

Increasing N while keeping Λ_H and r_π fixed one arrives at the following N-independent results:

$$\lim_{N\to\infty} S_{GB} = \left[\frac{f(m_Z/m_h)}{6\pi} + \frac{1}{12\pi}\log\frac{(4\pi\Lambda_H)^2}{m_h^2} + \frac{5}{72\pi}\right] + \frac{1}{12\pi}\log\frac{m_h^2}{m_{h,\rm ref}^2} + \frac{f(m_Z/m_{h,\rm ref}) - f(m_Z/m_h)}{6\pi}, \quad (19)$$

and

$$\lim_{N\to\infty} T_{GB} = -\frac{3}{16\pi\cos^2\theta_w}\log\frac{(4\pi\Lambda_H)^2}{m_h^2} + T_{\rm SM}^h(m_h) - T_{\rm SM}^h(m_{h,\rm ref}). \quad (20)$$

The corrections appear at $\mathcal{O}(N^2)$.

In Fig. 1 we plot S and T as a function of the number of colors for different values of Λ_H. Because the corrections are in $1/N^2$ the large N limit is approached quickly. We have also assumed $r_\pi \approx 1$, which is its natural order of magnitude and, in any event, can be partially reabsorbed in Λ_H. We compare the result with the experimental value of precision data in Fig. 2 for $\Lambda_H = 200$ GeV (blue curve) and

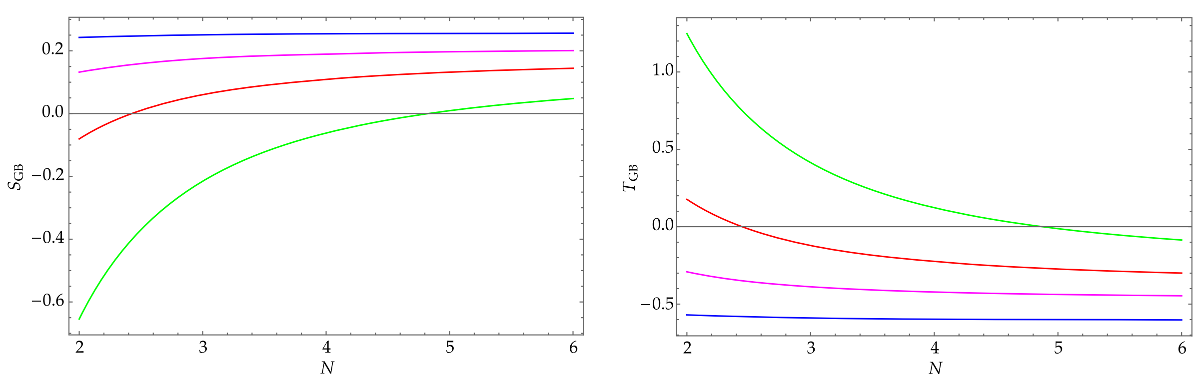

FIG. 1. We show the dependence on the number of underlying glueball-Higgs colors for the (left panel) S and (right panel) T for $\Lambda_H = 500$ (blue curve), 200 (magenta curve), 100 (red curve), and 50 GeV (green curve). The composite scale $4\pi\Lambda_H$ is always higher than the electroweak scale of 246 GeV, and $r_\pi \approx 1$ is further assumed.

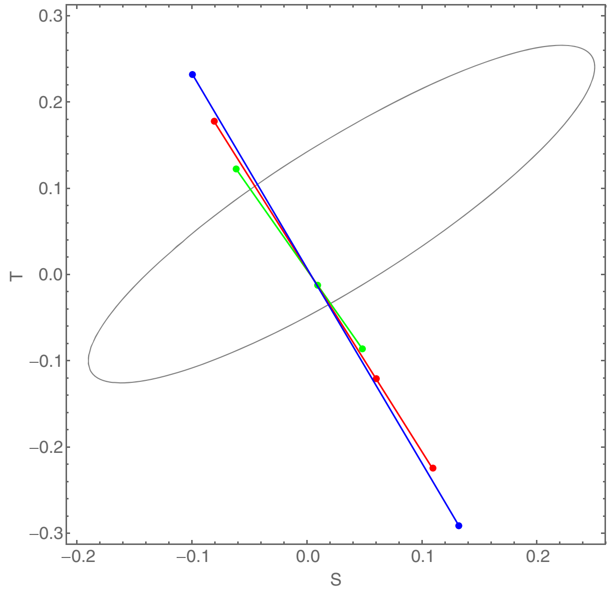

FIG. 2. Comparison with the precision electroweak constraints for the glueball-Higgs model for $\Lambda_H = 200$ GeV (blue curve) and $N = 1, 2$, $\Lambda_H = 100$ GeV (red curve) and $N = 2, 3, 4$, and $\Lambda_H = 50$ GeV (green curve) for $N = 4, 5, 6$. The leftmost point on each curve corresponds to the smallest N, and $r_\pi = 1$ is further assumed.

$N = 1, 2$. In red we have $\Lambda_H = 100$ GeV and $N = 2, 3, 4$. Finally we plot the 50 GeV (green curve) case for $N = 4, 5, 6$. Although we expect the composite scale to be around $4\pi\Lambda_H$, which even for this very low value of Λ_H is higher than the electroweak scale of 246 GeV, we need to consider higher order terms in the effective Lagrangian for a more consistent estimate for $\Lambda_H = 50$ GeV. The leftmost point on each curve corresponds to the smallest N. The experiments prefer smaller values of Λ_H with N in the range 2–4. Larger values of Λ_H require N to be away from the large N limit and therefore we cannot conclude on the viability of the $\Lambda_H = 200$ GeV case. Increasing further Λ_H it is clearly not preferred by precision observables. If, therefore, a glueball-Higgs model does describe the Higgs we expect soon new states to be discovered with masses in the range 600–1200 GeV.

We stress that by requiring them to be in agreement with precision measurements the couplings of the Higgs to the standard model gauge bosons are also close to the experimental values. This occurs because the product $N\Lambda_H$ is constrained to be near the electroweak scale.

For the dilaton-Higgs example we have

$$\kappa_1^{\rm Dilaton} = 2\frac{v}{N\Lambda_H}, \tag{21}$$

which corresponds to the results above but now with r_π exactly equal to one.

B. S and T for the large N dynamical Higgs

It is interesting to explore what happens for the large N dynamical Higgs. The main difference with respect to the previous case is that the electroweak scale and the dynamical Higgs scale are now identified. Among the possible underlying models that can lead to this kind of effective dynamics are time-honored examples such as minimal models of (near-conformal) technicolor [35–37],

$$\kappa_1^{TC} = 2r_\pi\sqrt{\frac{\bar{N}}{N}}, \tag{22}$$

yielding

$$S_{TC} = \left(1-\frac{\bar{N}}{N}r_\pi^2\right)\left[\frac{f(m_Z/m_h)}{6\pi}+\frac{1}{12\pi}\log\frac{(4\pi v)^2}{m_h^2}+\frac{5}{72\pi}\right] + \frac{1}{12\pi}\log\frac{m_h^2}{m_{h,\rm ref}^2}+\frac{f(m_Z/m_{h,\rm ref})-f(m_Z/m_h)}{6\pi}, \tag{23}$$

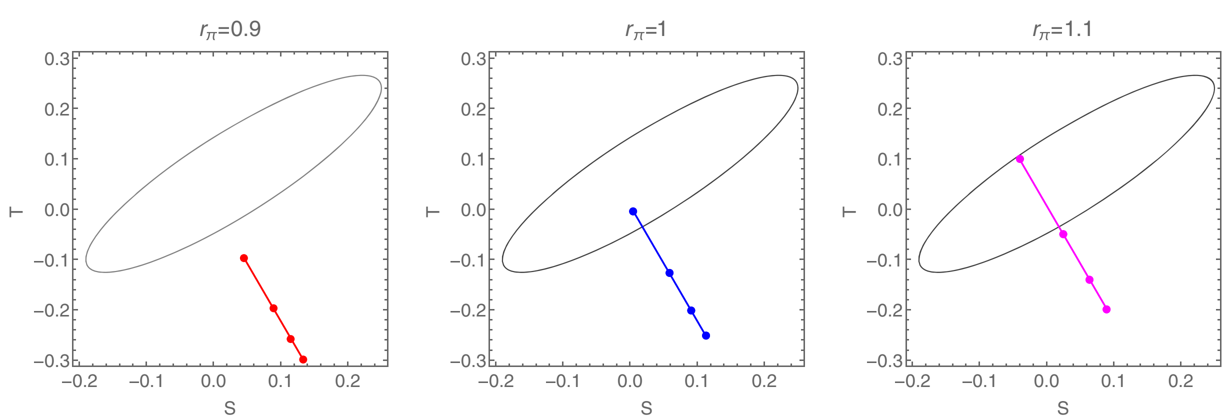

FIG. 3. Comparison with the precision electroweak constraints for the dynamical Higgs for $v = 246$ GeV and $N = 3, 4, 5, 6$. The leftmost point on each curve corresponds to the smallest N, and further assumed are $r_\pi = 0.9$ (left panel), $r_\pi = 1$ (center panel), $r_\pi \approx 1.1$ (right panel). The estimates are obtained using the effective Lagrangian in (10). We also use (9) and rename the left-hand side of the equation v to avoid introducing new symbols. Finally we have chosen the reference value $\bar{N} = 3$.

and

$$T_{TC} = -\frac{3}{16\pi\cos^2\theta_w}\left(1-\frac{\bar{N}}{N}r_\pi^2\right)\log\frac{(4\pi v)^2}{m_h^2} + T_{\mathrm{SM}}^h(m_h) - T_{\mathrm{SM}}^h(m_{h,\mathrm{ref}}). \quad (24)$$

Unlike the glueball-Higgs case we have at our disposal only the N dependence of the effective coupling, which goes to $1/N$ for the fundamental representation (chosen here) or to $1/N(N \pm 1)$ for two-index representations [29]. We show in Fig. 3 the comparison to the precision electroweak constraints for the dynamical Higgs for $v = 246$ GeV and $N = 3$, 4, 5, 6. To avoid the introduction of yet another symbol in this section the physical electroweak scale is labeled by v. The leftmost point on each curve corresponds to the smallest N, and further assumed are $r_\pi = 0.9$ (left panel), $r_\pi = 1$ (center panel), and $r_\pi \approx 1.1$ (right panel) It is clear from the results that it is possible to abide the electroweak precision constraints for a larger number of colors provided that r_π is larger than in the SM.

The computations show that finite $1/N$ corrections stemming from the nonstandard Higgs sector are relevant and cannot be neglected. These were not taken into account in [33]. The effects of heavier states such as these are included in the cutoff dependence of the counterterms.

V. CONCLUSIONS AND TOP CORRECTIONS

We introduced effective field theories and associated counting schemes to consistently describe the lightest massive large N stable composite scalar state emerging in any theory of composite dynamics. The framework allows for systematic investigations of composite dynamics featuring non-Goldstone (and Goldstone) scalars. As time-honored examples we discussed the lightest glueball state stemming from Yang-Mills theories. We further applied our effective approach to models of (near-conformal) dynamical electroweak symmetry breaking. In particular we considered the following three possibilities: the Higgs is the lightest glueball of a new composite theory; it is a large N scalar meson in models of dynamical Higgs such as technicolor; and finally that it is a large N pseudodilaton in the form of a conformal compensator. For each of these models, we provided the leading N corrections to the precision parameters.

We observe that it is straightforward to show that in this framework the top corrections to the glueball and dynamical Higgs mass can be reliably estimated in the large N limit by rescaling r_t in Eq. (4) of [38,39] by the opportune power of N, for each model, and simultaneously replacing the cutoff scale by either $4\pi\Lambda_H$ or $4\pi v$.

The results provide useful insights stemming from the large N dynamics of these models and can be viewed as the stepping stone for a consistent determination of quantum corrections at the effective Lagrangian level containing massive scalar states. The effective approach is directly applicable also to models of composite Goldstone Higgs dynamics [40,41] when including the first massive scalar state [39,42,43], as well as to investigate interesting flavor properties [44,45]. Finally holographic studies of the spectrum and large N properties of strongly coupled theories [46–48] can benefit from a model independent large N computation that can be performed with the effective theories constructed here.

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