Top Compositeness and Precision Unification

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The evolution of Standard Model gauge couplings is studied in a non-supersymmetric scenario in which the hierarchy problem is resolved by Higgs compositeness above the weak scale. It is argued that massiveness of the top quark combined with precision tests of the bottom quark imply that the right-handed top must also be composite. If, further, the Standard Model gauge symmetry is embedded into a simple subgroup of the unbroken composite-sector flavor symmetry, then precision coupling unification is shown to occur at $\sim 10^{15}$ GeV, to a degree comparable to supersymmetric unification.

The ambitious ideas of grand unification, and variants such as string unification and orbifold unification, are founded on the structure and successes of the Standard Model (SM). A central quantitative prediction is the evolution of SM gauge couplings from a single unified coupling, $\alpha_U$, at a unification scale, $M_U$. Schematically, in the minimal scenario,

$$\alpha_{i=1,2,3}(\mu) = \alpha_U + \text{SM} + M_U\text{-physics},$$

where the second term represents SM running, while the third represents model-dependent threshold effects from unification physics at a scale $\sim M_U$. Fortunately, it is natural for these $M_U$-scale effects to be much smaller than SM running, unless the $M_U$-sector is very large or has large non-degeneracies. Neglecting these effects allows one to test unification with just SM data.

There are several shortcomings: (i) The couplings do not meet very precisely. This does not falsify unification since $M_U$-effects may be unexpectedly large, but a more precise meeting without invoking such effects would have been much stronger circumstantial evidence for unification. Nevertheless, the results are intriguing. (ii) With the best fit, $M_U \sim 10^{14}$ GeV, so large that there is certainly no prospect of experimentally verifying any unified symmetry. (iii) The requisite SM extrapolation to such high $M_U$ results in a severe gauge hierarchy problem. (iv) This $M_U$ is still low enough that exchange of massive states can result in excessive proton decay. Such states can however be avoided in string or orbifold unifications.

By comparison, unification in the context of weak scale supersymmetry (SUSY) is a striking success. The coupling evolution is given schematically by

$$\alpha_i(\mu) = \alpha_U + \text{SM} + \text{superpartners} + M_U\text{-physics},$$

and again can be tested neglecting $M_U$-effects: (i) The superpartner-induced running yields a high precision meeting of couplings. The level of precision can be quantified by the postdiction $\delta_3 \equiv (\alpha_3^{\text{theory}} - \alpha_3^{\text{exp}})/\alpha_3^{\text{exp}} \sim 10\%$ at the scale $m_Z$. This size of $\delta_3$ can be naturally accounted for by threshold effects from $M_U$-physics. (ii) While the unification scale $M_U \sim 10^{16}$ GeV is still high, we rely less on directly seeing the unified gauge symmetry given the stronger circumstantial evidence. (iii) SUSY can solve the gauge hierarchy problem, so that two important issues are addressed simultaneously. (iv) $M_U$ is high enough to adequately suppress proton decay. String unification or orbifold unification are still attractive for solving the doublet-triplet splitting problem.

In this letter we pursue a very different scenario, namely that the hierarchy problem is solved by having the Higgs doublet be a composite of some new (non-supersymmetric) strong dynamics. While such dynamics is necessarily non-perturbative and theoretically challenging, there has been a recent revival of interest because of two extensions which allow one to understand weak scale symmetry breaking, precision tests and phenomenology, independent of many of the details of the strong sector. Little Higgs theory is one such extension, which we will not pursue here. Our work is motivated by (but not strongly reliant on) Refs. \cite{7, 8}, a realistic Randall-Sundrum (RS) extra-dimensional scenario with most or all of the SM fields in the bulk. Via the AdS/CFT correspondence, such a scenario is dual to a purely 4D composite Higgs scenario. The Kaluza-Klein excitations map to some low-lying hadrons at the compositeness scale, $\Lambda_{\text{comp}}$. The ratio of higher-dimensional curvature to the effective field theory cutoff maps to a new small parameter of the strong sector, with the help of which many weak-scale observables can be calculated independently of microscopic details of the strong dynamics.

An attractive feature of (the 4D dual of) this type of RS set-up is its simple extrapolation (at least for some important inclusive observables) to energies far above the weak scale. This leads to an elegant mechanism for gen-

\footnote{Often in the SUSY literature $\delta_3$ is evaluated at $M_U$ as being a few percent. This must be multiplied by $\alpha_3(m_Z)/\alpha_3(M_U) \sim 2.5$ in order to compare at $m_Z$, as we do.}

\footnote{We will not follow the technicolor approach, in which a Higgs scalar is effectively absent.}
The light SM fermions are taken to be elementary particles, weakly coupled to strong-sector operators. Running down to \( \Lambda_{\text{comp}} \), these operators induce small Yukawa couplings to the Higgs composite. The light SM fermions are elementary particles, weakly coupled to strong-sector operators. Running down to \( \Lambda_{\text{comp}} \), these operators induce small Yukawa couplings to the Higgs composite. Hierarchies arise naturally from the different scaling dimensions of different strong operators. The weak couplings to the strong sector also naturally suppress modifications of couplings to the W, Z and compositeness and flavor-changing effects in the light SM fermions, in accord with modern data.

The top quark is, however, a special case. Its Yukawa coupling to the Higgs composite is so large that either \( t_R \), \( t_L \), or both must effectively also be composite. However, precision data such as tests of \( Z \to b \bar{b} \) strongly suggest that \( b_L \), and hence \( t_L \) by electroweak symmetry, can have at most a small admixture of a TeV-scale composite. We deduce that \( t_R \) must be the composite.

With these broad motivations and expectations, we will show that under quite simple and plausible conditions an attractive scheme for precision unification emerges. We will first derive our central result to leading order (LO) in the couplings of elementary fields to the strong sector, and then consider subdominant corrections. Our discussion will be mostly from the 4D viewpoint of the strong sector, and then consider subdominant corrections. At one-SM-loop the \( H, t_R \) must be the composite.

Fortunately, the non-perturbative strong sector contributions to SM running cancel to one-SM-loop order in computing differential running, that is the running of \( (\alpha_i - \alpha_1) \) say, if the SM gauge group is embedded in a simple factor of \( G \) (such as \( SU(5) \) for example):

\[
\alpha_i(\mu) - \alpha_1(\mu) = \alpha_U + \text{SM} - \{ t_R, H \} + M_U-\text{physics}. \tag{4}
\]

This is all we need to check gauge coupling unification. We assume the simple embedding of the SM into \( G \) from now on.

Eq. 4 exhibits a remarkable twist in the unification paradigm. Instead of adding the running from physics beyond the SM, here compositeness instructs us to subtract the running due to some SM particles! Before checking unification we must explain why the light composites, \( H, t_R \), do not fill out complete representations of the global symmetry, \( G \). There are two distinct cases following from the possibility of spontaneous symmetry breaking in the strong sector at \( \Lambda_{\text{comp}} \), \( G \to K \). (Indeed the original, and still attractive, proposal for a composite Higgs is as a (pseudo-)Goldstone boson of this type of symmetry breaking, though this is not essential for the present paper.) The two cases are (a) the SM gauge group remains embedded in a simple factor of \( K \subset G \), or (b) it does not. In (b) there is no contradiction with \( H, t_R \) being the only light composites, and Eq. (4) applies. One finds that the subtractions certainly improve unification, but it is still not very precise. We will not study this case further here.

Here, we focus on (a), where \( H, t_R \) must be accompanied by other composites, filling out complete \( K \)-representations. Having extra (colored) scalar composites does not pose a robust problem, since the perturbations of the SM coupling to the strong sector can easily split the Higgs doublet from its \( K \)-partners, allowing the former to condense and be light while the latter do not condense and are massive enough to avoid present bounds. But the chiral fermionic \( K \)-partners of the \( t_R \) do pose a robust problem (and introduce SM anomalies). The only way to remove these unwanted states is to assume there exist exotic elementary fermions beyond the SM with couplings to the strong sector, which induce Dirac masses with the \( K \)-partners of the \( t_R \) below \( \Lambda_{\text{comp}} \). That is, the exotics must have SM quantum numbers which are charge-conjugate to the \( K \)-partners of the \( t_R \).

The elementary exotics also contribute to SM running above \( \Lambda_{\text{comp}} \),

\[
\alpha_i(\mu) = \alpha_U + \text{SM} - \{ t_R, H \} + \text{exotics} + \text{strong sector} + M_U-\text{physics}. \tag{5}
\]

We assume here that the exotic couplings to the strong sector are weak enough that their contributions to SM running are approximately undressed by strong-sector corrections. At one-SM-loop the differential running only depends on the fact that the exotics fill out a complete \( K \)-representation except for a missing \( t_R \).

\[
\alpha_i(\mu) - \alpha_1(\mu) = \alpha_U + \text{SM} - \{ t_R, t_R^c, H \} + M_U-\text{physics}. \tag{6}
\]

Neglecting the \( M_U \)-threshold, as usual, yields near perfect unification, with \( M_U \sim 10^{15} \text{ GeV} \). See Fig. 1.

Eq. (5) and Fig. 1 summarize our central quantitative result to LO in SM gauge couplings and zeroth order in the couplings of elementary fermions to the strong sector. We now discuss the subleading corrections. It is difficult to couple elementary fermions to strong sector operators at \( M_U \sim 10^{15} \text{ GeV} \) without the couplings being highly irrelevant in the IR, resulting in negligible
Yukawa couplings, unless the strong sector is strongly-coupled throughout the large hierarchy. This happens naturally when the strong sector is near an IR-attractive fixed point above $\Lambda_{\text{comp}}$. In this case, working to next-to-leading order (NLO), the gauge coupling running above $\Lambda_{\text{comp}}$ is given by

$$\frac{d}{d \ln \mu} \left( \frac{1}{\alpha_i} \right) = \frac{b_i}{2\pi} + \frac{B_{ij}}{2\pi} \frac{\alpha_j}{4\pi} + C_{i\alpha} \frac{\lambda_\alpha^2}{2\pi} \frac{1}{16\pi^2},$$

where the $b, B, C$ are constants and $\lambda_\alpha, \alpha = \text{exotic}, Q^2_\alpha \equiv (t_L, b_L)$, denote the largest couplings of the elementary fermions to the strong sector (resulting in the largest masses with composite fermions). We further decompose

$$b_i = b_i^{SM-} + b_i^{\text{exotic}} + b_i^{\text{strong}}$$
$$B_{ij} = B_{ij}^{SM-} + B_{ij}^{\text{exotic}} + B_{ij}^{\text{strong}},$$

where “SM−” refers to $SM - \{t_R, H\}$. Note that $b_i^{SM-}, b_i^{\text{exotic}}, B_{ij}^{SM-}, B_{ij}^{\text{exotic}}$ are just representation-theoretic factors. For concreteness we consider the SM gauge group embedded in SO(10) $\subset K$ in the usual way, with the $t_R$ being part of a composite 16 of SO(10), so that the elementary exotics $\equiv \{t_R\}$.

By contrast, $b_i^{\text{strong}}, B_{ij}^{\text{strong}}, C_{i\alpha}$ include unknown $O(1)$ strong interaction factors. We will treat $b_i^{\text{strong}}$ as an unknown ($i$-independent by the SM embedding into a simple factor of $G$), which can usefully be thought of as a crude measure of the SO(10)-charged content of the strong sector. A rough but reasonable expectation is that $b_i^{\text{strong}} \gtrsim b_i^{\text{comp}}$, where $b_i^{\text{comp}}$ is the LO renormalization group coefficient due to the light composites alone in the far IR. For a real scalar 10 and a Weyl fermion 16 of SO(10) ($\equiv H, t_R$), $b_i^{\text{comp}} = 1.5$. We will use crude estimates of the NLO coefficients, $B_{ij}^{\text{strong}} \sim 3 \cdot 3 \cdot b_i^{\text{strong}}$, $C_{i\alpha} \sim 3 \cdot b_i^{\text{strong}}$, as part of our theoretical error estimates. These estimates follow from the fact that the (non-perturbative) diagrams contributing to the NLO coefficients arise from diagrams contributing to $b_i^{\text{strong}}$ with insertions of intermediate elementary gauge bosons or fermions via SM gauge couplings or $\lambda_\alpha$. Such insertions can result in summation over QCD colors or weak isospins, giving rise to an extra factor of at most 3. Further, experience with perturbative gauge loops shows that they give an extra factor $\sim 3$ beyond the naive loop-counting parameter. This accounts for the second factor of 3 in $B_i^{\text{strong}}$.

Following (the AdS/CFT dual of) the scenario of Ref. [7], we assume that above $\Lambda_{\text{comp}}$ the couplings $\lambda_\alpha$ are slightly relevant, driving the theory away from the original fixed point (of the isolated strong sector) to a nearby fixed point. We can approximate $\lambda_\alpha(\mu)$ in the gauge coupling running by their new-fixed-point values, $\lambda_\alpha$. Integrating Eq. (7) down to $\mu \lesssim \Lambda_{\text{comp}}, m_{\text{exotic}}$, we find

$$\frac{1}{\alpha_i(\mu)} = \frac{1}{\alpha_U} - \frac{b_i}{2\pi} \ln \frac{M_U}{\mu} - \frac{C_{i\alpha}}{32\pi^3} \lambda_\alpha^2 \ln \frac{M_U}{\mu}$$
$$- \frac{B_{ij}}{4\pi b_j} \ln \left[ 1 + \lambda_j(\mu) \frac{b_j}{2\pi} \ln \frac{M_U}{\mu} \right] + \text{threshold corrections.}$$

Let us first discuss the thresholds at $M_U \sim 10^{15}$ GeV, $\Lambda_{\text{comp}} \sim$ few TeV, and $m_{\text{exotic}} \sim$ TeV. The detailed physics at $M_U$ is unknown, but the threshold effects associated with the strong sector can be subsumed into $\alpha_U, b_i^{\text{strong}}, C_{i\alpha}$. As in standard unification schemes, the minimal natural size of threshold effects associated with elementary fields is $\delta(1/\alpha_i) \sim O(1)/2\pi$. The expectation $\Lambda_{\text{comp}} \sim$ few TeV follows from the requirements of reasonable naturalness of the weak scale as well as passing electroweak precision tests. This was demonstrated in RS modelling using the extra-dimensional calculability [7, 8]. In the present scenario, we do not expect a useful extra-dimensional dual description, as we explain later, but the RS calculations of precision observable serve as plausible estimates, with at most $O(1)$ unknown correction factors. A central requirement however, is having an approximate “custodial isospin” symmetry of the strong sector to protect the electroweak $\rho$-parameter. Our choice of $SO(10)$ flavor symmetry ensures this, with custodial $SU(2)R$ as well as SM subgroups. As a consequence, one exotic has the gauge quantum numbers of $t_R$, is stable given a baryon number symmetry of the strong sector, and can serve as a dark matter candidate if its mass is $\sim$ few hundred GeV [22]. However, the exotic $SU(2)R$-partner of $t_R$ must have mass $\gtrsim 1.5$ TeV in order to avoid it forming too large a component of the observed bottom quark, in contradiction with precision tests. These considerations motivate $m_{\text{exotic}} \sim$ TeV, with mild SO(10)-violating splittings. A simple (but not the only) way for this to happen is for $SU(5)$ to be an exact flavor symmetry of the strong sector, with the remainder of the SO(10) symmetry being only approximate. Exact $SU(5)$ is all we need here.

FIG. 1: LO differential running of SM gauge couplings in the top/Higgs compositeness scenario (a).
In SUSY unification the one-loop superpartner threshold effects are generally significant because of large non-gauge-universal splittings in their spectrum induced by running from $M_U$, and also by the need to avoid existing search constraints and extreme fine-tuning. In the present scenario this does not happen because the $\Lambda_{\text{comp}}$ threshold approximately has the global $K$-symmetry of the strong sector, while the exotics and their Dirac partners also come in an almost $K$-symmetric form, with only the $t_R$ state missing. We will therefore probe our sensitivity to the associated threshold corrections by simply varying $\Lambda_{\text{comp}}$ from $3 - 5$ TeV, and insert a single (for simplicity) exotic threshold from $0.5 - 2$ TeV in running SM couplings measured at $m_Z$ up to $\Lambda_{\text{comp}}$.

Finally, we need to estimate the weakly-perturbed fixed-point couplings, $\lambda_{\alpha}$. $\lambda_{\alpha Q}^1$ is responsible for coupling the elementary $t_L$ to the strong sector, yielding a Yukawa coupling below $\Lambda_{\text{comp}}$ of $\lambda_{\alpha Q}^1$ times an $O(1)$ strong interaction factor. Thus we have $\lambda_{\alpha Q}^1 \sim 1$. $\lambda_{\text{exotic}}$ is responsible for generating a Dirac mass for the exotics with the excess fermion composites, $m_{\text{exotic}} \sim \lambda_{\text{exotic}} \Lambda_{\text{comp}} \sqrt{b^{\text{strong}}/4\pi}$. For $m_{\text{exotic}} \sim \text{(TeV)}$ we also need $\lambda_{\text{exotic}} \sim 1$. Thus, in our analysis $\lambda_{\alpha} \sim 1$.

In Fig. 2 we exhibit a simple and standard test of unification, given the high precision of electroweak data, namely using measured values of $\alpha_{1, 2}$ to postdict $\alpha_3(m_Z)$. We use separate bands to denote the variation in postdicted $\alpha_3(m_Z)$ coming from the above threshold ranges and from the theoretical error arising from our $B_{\text{strong}}^{\text{strong}}, C_{\text{IA}}$ bounds. Note that our central predictions are excellent, we do not need large corrections, our largest uncertainties just reflect the conservative bounds put on the $B_{\text{IA}}^{\text{strong}}, C$. The regime of controlled unification involves modest $b^{\text{strong}}$, not much larger than the size of the strong flavor group. Also, requiring that the SM gauge couplings do not have a Landau pole below $M_U$ implies $b^{\text{strong}} \lesssim 9$. This suggests that an AdS dual description, requiring a large ratio of strong colors ($\sim O(b^{\text{strong}})$) to flavors, will not be useful. This is why we have not pursued RS modelling in the present context.

Let us assess our scenario with the criteria used to discuss earlier unification scenarios: (i) The couplings meet very precisely, strong circumstantial evidence for this form of unification. The postdiction of $\alpha_3(m_Z)$ works to better than $\delta_3 \approx 15\%$ over a wide range of $b^{\text{strong}}$. Alternatively, postdiction of $\sin^2\theta_W$ using $\alpha_{\text{exp}}^{\text{expt}}$ leads to an error $\delta \sin^2\theta_W \approx 0.03 \delta_3$. This is quite comparable with the level of success of SUSY unification \cite{22}. (ii) With the best fit, $M_U \sim 10^{15}$ GeV, so large that experiments will not directly see the unification physics. However, there will be a striking signature of unification surviving to accessible energies \cite{18}, namely the strong sector resonances will fill out approximately degenerate multiplets of a unified flavor symmetry, $K$. Since their masses are expected to be in the few TeV range (based on naturalness), the cross-section for their single production at the LHC could be significant. These resonances will decay mostly into $t_R$ and Higgs (including longitudinal $W/Z$) due to the strong coupling involved, quite distinctly from other models such as SUSY. (iii) We have arrived at precision unification here by considering one of the simplest non-supersymmetric scenarios for solving the hierarchy problem of the SM. (iv) $M_U \sim 10^{15}$ GeV is still low enough that exchanges of $X, Y$ bosons at this scale can result in excessive proton decay. Even more importantly, composite states with the same quantum numbers can also mediate proton decay. The only known schemes in which both problems are solved are string or orbifold unification \cite{2, 3, 20}, so this is a requirement of our scenario. There could also be UV model-dependent states at $M_U$ contributing to proton decay. These effects can be suppressed by imposing a (gauged) baryon-number symmetry, compatible with orbifold unification, as long as it is broken somewhat below $M_U$. This symmetry should also be an accidental flavor symmetry of the strong dynamics, to extend the usual accidental baryon-number symmetry of the SM. A second reason for preferring string/orbifold unification is that it makes it simpler to understand the appearance of incomplete grand-unified fermion multiplets in the IR, such as our exotics. In orbifold unification the global strong-sector symmetry $G$ (or $K$) may even be the grand unified gauge group, surviving orbifold projections in this sector, but not in the elementary fermion/gauge-boson sectors. (v) As mentioned earlier, an attractive dark matter candidate emerges as a $K$-partner of $t_R$ \cite{22}.

The scenario in which the SM hierarchy problem is solved non-supersymmetrically with top/Higgs compositeness, or a RS dual depiction, is attractive from several phenomenological points of view. In this letter, we have studied one of the key features that has been taken as
strong evidence in favor of a supersymmetric solution to the hierarchy problem, namely precision gauge-coupling unification. We have found an equally striking (but very different) unification that follows rather minimally from top/Higgs compositeness. We hope to have shown that taking unification as a serious consideration, one must still keep an open mind as to how the hierarchy problem is resolved in Nature, supersymmetrically or non-supersymmetrically. This is not a passive state, extracting new physics from upcoming colliders is challenging and requires planning ahead.

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