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Minimal supersymmetric fat Higgs model

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We present a calculable supersymmetric theory of a composite "fat" Higgs boson. Electroweak symmetry is broken dynamically through a new gauge interaction that becomes strong at an intermediate scale. The Higgs boson mass can easily be 200–450 GeV along with the superpartner masses, solving the supersymmetric little hierarchy problem. We explicitly verify that the model is consistent with precision electroweak data without fine-tuning. Gauge coupling unification can be maintained despite the inherently strong dynamics involved in electroweak symmetry breaking. Supersymmetrizing the standard model therefore does not imply a light Higgs boson mass, contrary to the lore in the literature. The Higgs sector of the minimal fat Higgs model has a mass spectrum that is distinctly different from that of the minimal supersymmetric standard model.

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I. INTRODUCTION

The mechanism and dynamics of electroweak symmetry breaking [1] is one of the greatest mysteries in particle physics. The standard model (SM) provides an extremely successful parameterization of electroweak symmetry breaking through the introduction of a single Higgs doublet with a negative mass squared. Furthermore, fermion masses, mixings and *CP* violation are accommodated in a way that is consistent with experimental observations and constraints from flavor-changing neutral current processes. Unfortunately, the negative mass squared is just an input to the theory without a deeper understanding. Moreover, symmetry breaking by relevant operators is fraught with an extreme ultraviolet (UV) sensitivity: the Higgs boson mass squared is quadratically sensitive to new physics. Thus the origin of electroweak symmetry breaking in the SM is impossible to understand without a complete UV theory.

It is well known that supersymmetry removes the extreme quadratic sensitivity to the UV and can explain electroweak symmetry breaking with clear predictions for experiments. The minimal supersymmetric standard model (MSSM), in particular, predicts a light Higgs boson, $m_h \lesssim 130$ GeV, for top-squark masses less than about 1 TeV. Superpartners, especially charginos and top squarks, are also expected to be light because their masses feed into the Higgs boson mass parameters via the renormalization group [2]. The experimental lower bounds on the mass of the lightest Higgs boson and the masses of the superpartners are increasingly becoming a quantifiable concern. Even though the MSSM is far from being excluded, it already relies on fine-tuning at the level of a few percent: we call this the "supersymmetric little hierarchy problem."

The simplest way around this problem is to raise the Higgs boson mass. In the MSSM the Higgs boson mass could be raised by increasing the top squark masses, but only with an unnaturally drastic increase in fine-tuning. A less fine-tuned approach is to invoke physics beyond the MSSM that provides additional contributions to the quartic coupling. The simplest idea of this class is to add an extra singlet with a new (undetermined) Yukawa interaction with the Higgs fields. This is sufficient to raise the Higgs boson mass from 130 GeV up to about 150 GeV with fine-tuning comparable to that in the MSSM. But any further upward push on the Higgs boson mass causes the new Yukawa coupling to blow up at a scale below the gauge coupling unification scale.

In fact, incorporating a significantly heavier Higgs boson mass into beyond-the-MSSM models invariably leads to some type of strong coupling behavior, such as a Landau pole in the next-to-minimal supersymmetric standard model (NMSSM), e.g. [3], or simply a low UV cutoff, e.g. [4]. This has usually been viewed negatively because it ruins the UV completeness of supersymmetric theories, and eliminates connections to attractive theoretical ideas such as grand unification and string theory. The usual lore is that strong coupling should be avoided, yielding an upper bound on the mass of the Higgs boson.

We propose a radical revision of the usual lore and allow the Higgs sector to become strongly coupled at an intermediate scale. At the scale of strong coupling the Higgs fields reveal their composite nature and are in fact mesons of a confining theory in the UV. This theory is renormalizable, asymptotically free, and thus UV complete. The IR and UV dynamics are completely under control thanks to the improved knowledge of strongly coupled supersymmetric gauge theories [5] which is reliable even when soft supersymmetry breaking is added as a small perturbation on the strong dynamics [6,7]. We draw significant inspiration from recent proposals to fuse supersymmetry with technicolor [8,9], and indeed in our model electroweak symmetry is bro-

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ken dynamically. However, we have physical (composite) Higgs fields in the low energy effective theory with no *a priori* restriction on the scale of strong coupling, reminiscent of the older non-supersymmetric composite Higgs models [10] (but without the associated fine-tuning problems).

The outline of this paper is as follows. We first discuss the supersymmetric little hierarchy problem in Sec. II, and emphasize that a heavier Higgs boson mass easily solves this problem. We then construct our supersymmetric composite Higgs theory in Sec. III. Our basic framework is a three-flavor $SU(2)_H$ theory that *s*-confines, resulting in a low energy effective Lagrangian containing a dynamically generated superpotential of composite mesons. By introducing a mass for one flavor below the compositeness scale, we show that the mesons acquire expectation values that break electroweak symmetry at a scale that is tunable through this mass parameter. In Sec. IV we demonstrate how the various energy scales can be naturally obtained from the supersymmetry breaking, and in Sec. V we show how fermion masses and mixings can be incorporated. We calculate the scalar spectrum in Sec. VI. In Sec. VII we briefly comment on the new phenomenology of this model, emphasizing the unusual scalar spectrum. Section VIII explains how gauge coupling unification can be preserved. Finally, we conclude with a discussion in Sec. IX.

II. SUPERSYMMETRIC LITTLE HIERARCHY PROBLEM

In this section we define the supersymmetric little hierarchy problem and propose a simple but unconventional way out. The problem is that the conventional supersymmetric theories are increasingly fine-tuned since the Higgs boson and/or charginos have not yet been discovered. We point out that a composite Higgs boson will solve this problem easily once a suitable UV completion is found.

The MSSM provides a simple way to understand electroweak symmetry breaking (see, e.g., [11] for a review). In the supersymmetric limit, electroweak symmetry is not broken. Therefore, electroweak breaking is solely due to the soft supersymmetry breaking effects. It arises from the renormalization of the up-type Higgs boson soft mass squared that is driven negative by the top Yukawa coupling. At the one-loop approximation, one finds

$$\Delta m_{H_u}^2 \sim -12\frac{h_t^2}{16\pi^2} m_{\tilde{t}}^2 \log\frac{M_{\rm UV}}{\mu_{\rm IR}}, \qquad (1)$$

where $M_{\rm UV}$ ($\mu_{\rm IR}$) is the UV (IR) cutoff and h_t is the top Yukawa coupling. Even with the universal boundary condition $m_{\tilde{t}} = m_{H_u}$, it is easy to see that a large logarithm between the weak scale and, say, the grand unification theory (GUT) scale makes $m_{H_u}^2$ negative. Assuming that the supersymmetric mass μ for the Higgs doublets is smaller in magnitude than m_{H_u}, electroweak symmetry is broken. This so-called radiative breaking of the electroweak symmetry is a very nice feature of the MSSM.

However, the phenomenological situation is forcing some degree of fine-tuning on the MSSM in the following fashion. First of all, the Higgs quartic coupling is given only by D terms that are determined by the electroweak gauge couplings

$$V_D = \frac{g^2+g'^2}{8}(|H_u|^2-|H_d|^2)^2. \qquad (2)$$

This implies that the natural scale for the Higgs boson mass is m_Z, and indeed there is a well-known tree-level upper limit on the lightest Higgs boson mass that is precisely m_Z. The only way to increase the Higgs boson mass is by using the $O(h_t^4)$ radiative correction to the Higgs boson quartic coupling. The approximate formula valid for a moderate $\tan\beta$ is

$$m_{h^0}^2 \simeq m_Z^2 + \frac{3}{4\pi^2} h_t^4 v^2 \log\frac{m_{\tilde{t}_1} m_{\tilde{t}_2}}{m_t^2}. \qquad (3)$$

Here, $v = 174$ GeV. Because the Higgs boson has not been found up to 115 GeV, this implies $m_{\tilde{t}} \gtrsim 500$ GeV. On the other hand, the minimization of the scalar potential leads to

$$\frac{1}{2}m_Z^2 = -\mu^2 - \frac{m_{H_u}^2 \tan^2\beta - m_{H_d}^2}{\tan^2\beta - 1} \simeq -\mu^2 - m_{H_u}^2, \qquad (4)$$

again for moderate $\tan\beta$. Therefore we need to fine-tune the bare $m_{H_u}^2$ and/or μ against the radiative correction in Eq. (1) at the level of

$$\frac{|\Delta m_{H_u}^2|}{m_Z^2/2} \sim 4.8\left(\frac{m_{\tilde{t}}}{500\ \text{GeV}}\right)^2 \log\frac{M_{\rm UV}}{\mu_{\rm IR}}. \qquad (5)$$

Even for a low UV scale of $M_{\rm UV} = 100$ TeV, this already requires a fine-tuning of 3%.

In addition, the null results from searches for charginos at the CERN e^+e^- collider LEP-II give a lower bound $M_2 \gtrsim 100$ GeV. Assuming a GUT relation among the gaugino masses, this implies $M_3 \gtrsim 350$ GeV. Because M_3 feeds into $m_{\tilde{t}}$ through renormalization group evolution, this then feeds into $m_{H_u}^2$, aggravating the situation. Moreover, the MSSM potential is rather delicate due to the possible instability along the D-flat direction $H_u = H_d$.

The situation would clearly be better if the tree-level Higgs boson mass could be raised above the LEP bound. Modifying Eq. (4), however, necessarily involves additional contributions to the Higgs potential that are not related to the SM gauge couplings. Furthermore, reducing the need for top (s)quark contributions to electroweak symmetry breaking and the Higgs boson mass, Eqs. (1) and (3), respectively, may help reduce the fine-tuning required. We will see that the fat Higgs model we propose in this paper achieves both of these aims.

The simplest extension of the MSSM that raises the tree-level Higgs boson mass is the NMSSM. In the NMSSM the μ term is replaced by the superpotential

$$W = \lambda N H_u H_d - \frac{k}{3} N^3 \tag{6}$$

where N is neutral under the SM and λ is undetermined. The Higgs boson quartic coupling therefore depends on λ as well as the gauge couplings, potentially allowing for a much higher Higgs boson mass. Increasing the Higgs mass requires a large λ. This coupling renormalizes *upward* with increasing energy, eventually encountering a Landau pole. At this scale the perturbative description breaks down and the theory is no longer UV complete. To avoid this problem, it is customary to impose the requirement that all coupling constants, and in particular λ, remain perturbative up to the gauge coupling unification scale. This places an upper bound on λ leading to an upper bound on the lightest Higgs boson mass of about 150 GeV [3,12–15]. Adding more matter fields can relax the bound somewhat, but not much [16–18]. Even for extensions of the NMSSM with Higgs fields in other representations this bound is relaxed to at most $m_h \lesssim 200$ GeV [19]. This is the basis for the lore that the lightest Higgs boson mass cannot be much higher than in the MSSM.[1]

This is the supersymmetric little hierarchy problem we intend to solve. We are seeking a theory where the Higgs boson mass is allowed to be much larger than m_Z at tree level. It is clear that the heart of the problem in the MSSM is that the quartic coupling is determined by the gauge couplings plus radiative corrections. This is the ultimate source of the tension between the top squark masses and the lightest Higgs boson mass. If the tree-level Higgs boson mass can be higher there is no need to rely on the radiative corrections from the top-quark–top-squark sector and therefore top squarks can be light. In fact, if all superpartners are around 200–450 GeV, the natural scale for the Higgs boson soft mass is also around the same scale, and there is no fine-tuning.

In this paper we employ exact results in supersymmetric gauge theories to UV complete a strongly coupled Higgs sector. In our model the Higgs fields are composite bound states of fundamental fields charged under a new, strongly coupled gauge theory. The supersymmetric strong dynamics drive electroweak symmetry breaking. The effective theory of the Higgs boson composites is a variant of the NMSSM with an arbitrarily strong quartic coupling. Furthermore, unlike the MSSM, the modified potential has no flat directions that may cause an instability.

III. THE FAT HIGGS MODEL

First we describe the dynamics that leads to electroweak symmetry breaking with composite "fat" Higgs fields. The model is an $N=1$ supersymmetric $SU(2)$ gauge theory with six doublets, $T^1, \ldots, T^6$. They carry the quantum numbers given in Table I.

TABLE I. Field content under $SU(2)_L \times SU(2)_H$ gauge and $SU(2)_R \times SU(2)_g \times U(1)_R$ global symmetries. The $U(1)_Y$ subgroup of $SU(2)_R$ is gauged.

Superfields	$SU(2)_L$	$SU(2)_H$	$SU(2)_R$	$SU(2)_g$	$U(1)_R$
$(T^1,T^2)\equiv T$	**2**	**2**	**1**	**1**	0
(T^3,T^4)	**1**	**2**	**2**	**1**	0
(T^5,T^6)	**1**	**2**	**1**	**2**	1
P	**2**	**1**	**1**	**2**	1
Q	**1**	**1**	**2**	**2**	1
S	**1**	**1**	**1**	**1**	2
S'	**1**	**1**	**1**	**1**	2

The tree-level superpotential consists of several terms,

$$W = W_1 + W_2 + W_3, \tag{7}$$

where

$$W_1 = ySnumT^1T^2 + yS'T^3T^4, \tag{8}$$

$$W_2 = -mT^5T^6, \tag{9}$$

$$W_3 = y(T^1,T^2)P\begin{pmatrix} T^5 \\ T^6 \end{pmatrix} + y(T^3,T^4)Q\begin{pmatrix} T^5 \\ T^6 \end{pmatrix}. \tag{10}$$

The singlet fields S and S' in W_1 are introduced to ensure that electroweak symmetry is broken. They are really optional since the condensate may well form along an electroweak breaking direction without forcing it through the F-flatness conditions of S and S'. W_2 is simply a mass term for the fifth and sixth doublets. The mass parameter m controls the separation between the electroweak breaking scale and the compositeness scale, as we will show. Finally, W_3 contains the fields P and Q which are 2×2 matrices that transform as doublets under $SU(2)_L$ or $SU(2)_R$ and also a global $SU(2)_g$. They are present simply to marry off certain "spectator" composite fields with a mass of order the compositeness scale. W_3 is optional, since these spectator composite fields also acquire electroweak symmetry breaking masses. But the addition of the P and Q fields with the above superpotential has the benefit of minimizing the field content of the low energy effective theory, which we call the minimal supersymmetric fat Higgs model.

Note also that the overall superpotential is natural if we assign non-anomalous $U(1)_R$ charges as shown in Table I. The global symmetries still allow for linear terms in S and S' in the superpotential that could be forbidden by an additional Z_3 symmetry. This also prevents tadpole diagrams for the singlets that could destabilize the weak scale [22]. Mass terms for the first four doublets, if present, could be eliminated by shifting the fields S and S'. Our model is not sensitive to the precise Yukawa couplings in the superpotential, but for simplicity we take a common y that is assumed to have the bare value $y_0 \sim O(1)$.

[1]In [20] a heavier Higgs boson mass was claimed possible if a top squark bound state condenses due to a strong trilinear coupling. In [21] a new moderately strong gauge interaction was used to enhance the Higgs boson quartic coupling, assuming a large supersymmetry breaking of about 7 TeV in a part of the theory.

$SU(2)_H$ becomes strong at a scale Λ_H. The theory has six doublets, so below Λ_H it is described by meson composites $M_{ij}=T^iT^j$ ($i,j=1,\ldots,6$) with a dynamically generated superpotential PfM/Λ^3. Together with the tree-level terms,

$$W_{\rm eff}=\frac{{\rm Pf}M}{\Lambda^3}-mM_{56}+ySM_{12}+yS'M_{34}$$
$$+yP^{k,\alpha}M_{k,\alpha+4}+yQ^{l,\alpha}M_{l+2,\alpha+4} \qquad (11)$$

where $k=1,2$ is an $SU(2)_L$ index contracted with P; $l=1,2$ is an $SU(2)_R$ index contracted with Q; and $\alpha=1,2$ is an $SU(2)_g$ index. In terms of the canonically normalized fields this becomes

$$W_{dyn}=\lambda({\rm Pf}M-v_0^2M_{56})+m_{\rm spect}(SM_{12}+S'M_{34}$$
$$+P^{k,\alpha}M_{k,\alpha+4}+Q^{l,\alpha}M_{l+2,\alpha+4}), \qquad (12)$$

where naive dimensional analysis (NDA) [23] suggests[2]

$$v_0^2\sim\frac{m\Lambda_H}{(4\pi)^2}, \qquad (13)$$

$$m_{\rm spect}\sim y\frac{\Lambda_H}{4\pi}, \qquad (14)$$

$$\lambda(\Lambda_H)\sim 4\pi. \qquad (15)$$

The crucial observation is that the scale of electroweak symmetry breaking, v_0, is generated dynamically and is controlled by the value of the supersymmetric mass m.

It is useful to change the notation for the meson matrix to make the role of different components clear:

$$N=M_{56}, \quad \begin{pmatrix} H_u^+ \\ H_u^0 \end{pmatrix}=\begin{pmatrix} M_{13} \\ M_{23} \end{pmatrix}, \quad \begin{pmatrix} H_d^0 \\ H_d^- \end{pmatrix}=\begin{pmatrix} M_{14} \\ M_{24} \end{pmatrix}, \qquad (16)$$

and all the other components of the meson matrix decouple near Λ_H. These Higgs fields have the dynamically generated superpotential

$$W=\lambda N(H_dH_u-v_0^2). \qquad (17)$$

The H_d and H_u doublets play the role of the MSSM Higgs doublets. As advertised, this superpotential forces electroweak symmetry breaking without relying on supersymmetry breaking effects. This superpotential was also studied in Ref. [24].

The strong coupling λ rapidly renormalizes to smaller values as the energy scale is reduced. We can estimate this running by neglecting corrections from gauge couplings and the top coupling. The solution to the one-loop renormalization group equation is[3]

$$\lambda^2(t)=\frac{2\pi^2}{2\pi^2\lambda^{-2}(0)+t}, \qquad (18)$$

where $t\equiv\log(\Lambda_H/\mu)$ *increases* toward the infrared. Using the NDA estimate $\lambda(0)\sim 4\pi$, we find, for example, $\lambda(4.5)\simeq 2$, practically independent of the initial value. If m is well below Λ_H, condensation occurs at the scale $4\pi v_0\sim(m\Lambda_H)^{1/2}\ll\Lambda_H$ where the theory is weakly coupled and therefore calculable.

This is one of our important results. In the infrared the theory contains Higgs states with a weakly coupled, renormalizable superpotential described by just two parameters, λ and v_0. This is a rather nontrivial result that depends on the specific choice of an $SU(2)_H$ gauge theory with three flavors. Other choices are not so interesting. For example, precision electroweak constraints tend to severely restrict new gauge interactions beyond the SM at the TeV scale, and thus theories with a dual magnetic description are not good candidates. For an $SU(N_c)$ theory there is a dynamically generated superpotential when $N_f=N_c+1$, and its renormalizability requires $N_f\leqslant 3$. This means that an $SU(2)$ gauge theory with three flavors is the unique choice for this purpose.

IV. SCALES

Phenomenologically, the scale of supersymmetry breaking soft masses must be near the electroweak scale, $\lambda v_0\sim m_{\rm SUSY}$, because much larger supersymmetry (SUSY) breaking would lead to fine-tuning, whereas a much smaller SUSY-breaking scale would have already been observed. Using the parameters of the UV theory, this implies $m\Lambda_H\sim(4\pi m_{\rm SUSY})^2$. This coincidence of scales is reminiscent of the μ problem in the MSSM. Here we show that this can be naturally obtained by a combination of the seesaw mechanism [25] and the Giudice-Masiero mechanism [26]. This requires conventional gravity-mediated supersymmetry breaking, which we assume for the discussion in this section.

The simplest way to relate Λ_H to other scales is to employ a superconformal theory where the gauge coupling remains constant over many decades in energy. We introduce two extra doublets T^7 and T^8 to $SU(2)_H$ for this purpose. We assume that T^7 and T^8 transform as a vector-like pair under other symmetries so that a supersymmetric mass term can be added to the superpotential,

$$W=m'T^7T^8. \qquad (19)$$

This theory with $N_c=2$ and $N_f=4$ is in the superconformal window [5]. At some scale Λ_4 the $SU(2)_H$ gauge coupling becomes strong and remains strong all the way down to m',

[2]Here, the parameters are defined at the scale Λ_H, and hence are not the bare ones. As we will see in the next section, $m\sim 4\pi m_0$, $y\sim 4\pi y_0$, and $\Lambda_H\sim m'\sim 4\pi m_0'$, due to the superconformal dynamics.

[3]This formula assumes that the spectators decouple at scale Λ_H. This is justified in the next section where the superconformal dynamics enhances y to $y\sim 4\pi y_0\sim 4\pi$.

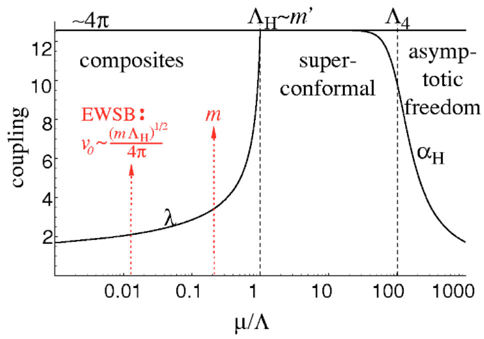

FIG. 1. The renormalization of the couplings in our fat Higgs model. The model becomes strong and nearly conformal at the scale Λ_4, where α_H nears 4π. The conformal invariance is broken by the mass of the extra doublet, m', which makes the theory confine at $\Lambda_H \sim m'$. Below this scale the effective theory description becomes one of meson composites with a coupling λ that quickly renormalizes down to $O(1)$. When $4\pi v_0 \ll \Lambda_H$ the mesons condense at weak coupling and the theory is calculable.

the supersymmetric vector-like mass of the extra doublets. At the scale m' the conformal symmetry is broken and $T^{7,8}$ may be integrated out. Below this scale the theory confines and is effectively the three flavor model discussed in the previous section. We therefore identify the strong coupling scale Λ_H with m'. The renormalization group evolution of the couplings is schematically shown in Fig. 1.

In addition to determining the scale Λ_H, the conformal dynamics generate large anomalous dimensions which have the effect of enhancing the couplings of the T fields, and therefore also the couplings of the composite Higgs fields. The structure of the superconformal algebra determines the anomalous dimensions exactly in terms of the anomaly-free R charges. Running from the strong scale Λ_4 down to the scale of conformal breaking Λ_H, the wave function of the T's is suppressed as

$$Z \sim \left(\frac{\Lambda_H}{\Lambda_4}\right)^{\gamma_*} \tag{20}$$

where $\gamma_* = 1/2$ is the anomalous dimension. Once the fields are canonically normalized this leads to an *enhancement* of their couplings. For example, the effective mass m' gets enhanced by a factor of

$$\left(\frac{\Lambda_4}{\Lambda_H}\right)^{1/2}. \tag{21}$$

In the low energy theory, any operator that involves one Higgs field, such as the top Yukawa, will be enhanced by a similar factor. Because the superconformal dynamics is likely to be upset by other strong couplings, the largest enhancement factor we consider is 4π.

The next task is to determine how m of the right size can be generated. First, it is assumed that the heavier vector-like mass m' is unrelated to supersymmetry breaking and therefore arbitrary. The scale for m' is presumably set by other flavor symmetries, akin to the right-handed neutrino mass which is protected by lepton number. However, the symmetries may conspire to forbid a vector-like mass m for the third flavor, analogous to the left-handed neutrino mass in the neutrino mass matrix. For example, consider a simple $U(1)$ flavor symmetry of charge $+1$ (-1) for the third (fourth) flavor. The symmetry is broken by an order parameter of charge $+2$. Then m' is allowed in the superpotential while m is not. Nevertheless, mixing between the third and fourth flavors is allowed by the symmetries and originates from the supersymmetry breaking due to the Giudice-Masiero mechanism. Therefore, the form of the mass matrix for these flavors becomes

$$\begin{pmatrix} 0 & m_{\mathrm{SUSY}} \\ m_{\mathrm{SUSY}} & m' \end{pmatrix}. \tag{22}$$

The light eigenvalue is given by $m = m_{\mathrm{SUSY}}^2/m'$. After the conformal dynamics enhances both m and m', we naturally obtain $mm' \sim (4\pi m_{\mathrm{SUSY}})^2$ as desired.

V. FERMION MASSES

In order to incorporate fermion masses, we follow [9] by adding four additional chiral multiplets that are singlets under $SU(2)_H$ but have the same quantum numbers as the Higgs doublets H_u and H_d in the MSSM,

$$\varphi_u, \overline{\varphi}_d\left(\mathbf{1},\mathbf{2},+\frac{1}{2}\right), \quad \varphi_d, \overline{\varphi}_u\left(\mathbf{1},\mathbf{2},-\frac{1}{2}\right). \tag{23}$$

They have the superpotential

$$\begin{aligned} W_f = {} & M_f(\varphi_u\overline{\varphi}_u + \overline{\varphi}_d\varphi_d) + \overline{\varphi}_d(TT^4) + \overline{\varphi}_u(TT^3) \\ & + h_u^{ij}Q_iu_j\varphi_u + h_d^{ij}Q_id_j\varphi_d + h_e^{ij}L_ie_j\varphi_d, \end{aligned} \tag{24}$$

where M_f is the mass of φ and $\overline{\varphi}$. The only flavor-violating couplings are the Yukawa couplings h_u^{ij}, h_d^{ij}, h_e^{ij}. We assume $M_f \sim m' \sim \Lambda_H$, possibly due to the same flavor symmetries that control the size of m'.

Between Λ_4 and $\Lambda_H \simeq m'$ the superconformal dynamics enhances the Yukawa couplings by $(\Lambda_4/\Lambda_H)^{1/2} \sim 4\pi$, as described in the previous section. After the φ's are integrated out, the effective dimension-5 superpotential is

$$W_f = \frac{4\pi}{M_f}[h_u^{ij}Q_iu_j(TT^3) + h_d^{ij}Q_id_j(TT^4) + h_e^{ij}L_ie_j(TT^4)]. \tag{25}$$

Below the compositeness scale Λ_H, NDA specifies the replacement $(TT^3) \to \Lambda_H H_u/4\pi$, $(TT^4) \to \Lambda_H H_d/4\pi$. Using $M_f \sim \Lambda_H$, the superpotential becomes

$$W_f = h_u^{ij}Q_iu_jH_u + h_d^{ij}Q_id_jH_d + h_e^{ij}L_ie_jH_d. \tag{26}$$

One may wonder if the Yukawa couplings are suppressed in the low-energy theory due to the wave function renormalization of the Higgs fields due to the strong coupling λ. Again using the one-loop renormalization group equation for simplicity, we find

$$h(t)=h(0)\left(\frac{\lambda(t)}{\lambda(0)}\right)^{1/4}. \qquad (27)$$

For $\lambda(0)\sim 4\pi$ and $\lambda(t)\sim 2$, we find that the suppression is only 60%. Therefore the mechanism presented above does yield sufficiently large Yukawa couplings.

VI. HIGGS BOSON MASS SPECTRUM

In this section the mass spectrum of the model is calculated. The supersymmetric part of the Higgs potential is

$$V_{\mathrm{SUSY}}=\lambda^2|H_dH_u-v_0^2|^2+\lambda^2|N|^2(|H_u|^2+|H_d|^2), \quad (28)$$

together with the D-term contributions that are familiar from the MSSM,

$$V_D=\frac{g^2}{8}(H_u^\dagger\vec{\tau}H_u+H_d^\dagger\vec{\tau}H_d)^2+\frac{g'^2}{8}(|H_u|^2-|H_d|^2)^2. \qquad (29)$$

Unlike the MSSM, electroweak symmetry breaking is caused by the confining dynamics even in the absence of supersymmetry breaking. Nevertheless, the potential also contains soft supersymmetry breaking terms

$$\begin{aligned}V_{\mathrm{soft}}&=m_1^2|H_d|^2+m_2^2|H_u|^2+m_0^2|N|^2\\&\quad+(A\lambda NH_dH_u-C\lambda v_0^2N+\mathrm{H.c.})\end{aligned} \qquad (30)$$

where $m_1,m_2,m_0,A,C\sim m_{\mathrm{SUSY}}$.

It is instructive to first look at a simple case where $m_1=m_2=m_0$, $A=C=0$. In this case, we can define the "standard model–like Higgs boson" $H=(H_u^0+H_d^0)/\sqrt{2}$ whose potential is simply

$$\begin{aligned}V&=\frac{1}{4}\lambda^2|H^2-2v_0^2|^2+m_0^2|H|^2\\&=\frac{1}{4}\lambda^2|H^2|^2-(\lambda^2v_0^2-m_0^2)|H|^2+\mathrm{const.}\end{aligned} \qquad (31)$$

This is no different from the potential in the minimal standard model. It is clear that electroweak symmetry is broken so long as $\lambda^2v_0^2>m_0^2$. The vacuum expectation value is $v=\langle H\rangle=\sqrt{2(v_0^2-m_0^2/\lambda^2)}$, and the mass of the Higgs boson is λv. There is an exact custodial $SU(2)$ symmetry in the Higgs sector.

For a more general set of parameters it becomes difficult to solve for the vacuum analytically. We take advantage of the (small) hierarchy

$$\lambda v_0\sim m_{1,2}\gg gv_0,g'v_0, \qquad (32)$$

which allows us to drop the MSSM D terms. If m_1/m_2 is very large, the quartic term in the potential is dominated by the D term and cannot be ignored. Our solution for the vacuum state applies for a moderate ratio m_1/m_2 where both Higgs boson masses $m_{1,2}$ are much larger than m_Z.

The fact that the Higgs boson quartic coupling of the MSSM is negligible compared to that coming from the strong dynamics illustrates that our model manifestly solves the supersymmetric little hierarchy problem, as anticipated in Sec. II. For simplicity we also set $A=C=0$ for most of the discussions below. With these approximations, the ground state is

$$H_u^0=v_0\sqrt{1-\frac{m_1m_2}{\lambda^2v_0^2}}\sqrt{\frac{m_1}{m_2}}, \qquad (33)$$

$$H_d^0=v_0\sqrt{1-\frac{m_1m_2}{\lambda^2v_0^2}}\sqrt{\frac{m_2}{m_1}}, \qquad (34)$$

$$N=0 \qquad (35)$$

up to corrections of order $m_Z^2/m_{1,2}^2$. To leading order in A and C we find that N no longer vanishes,

$$N=\frac{m_1m_2[(-A+C)\lambda^2v_0^2+Am_1m_2]}{\lambda[(\lambda^2v_0^2-m_1m_2)(m_1^2+m_2^2)+m_1m_2m_0^2]}, \qquad (36)$$

while shifts to H_u and H_d are only $O(A^2,AC,C^2)$. This demonstrates that a μ term is naturally generated, giving a mass of order m_{SUSY} to the Higgsinos.

Our vacuum solution was obtained by assuming that $m_{0,1,2}^2>0$, and thus arises from dynamical (as opposed to radiative) breaking of electroweak symmetry. Nevertheless, we expect a stable vacuum with electroweak symmetry breaking even if some or all of the (mass)^2 are negative because our potential Eqs. (28)–(31) is bounded from below, unlike in the MSSM where there is a possible instability along the D-flat direction. We leave such cases for a future study.

It is rather convenient to define

$$m_s=\sqrt{m_1m_2}, \qquad (37)$$

$$\tan\beta\equiv\frac{\langle H_u^0\rangle}{\langle H_d^0\rangle}=\frac{m_1}{m_2}, \qquad (38)$$

and then the electroweak breaking scale $v\simeq 174$ GeV is fixed in terms of the parameters of the model,

$$v^2=2\frac{\lambda^2v_0^2-m_s^2}{\lambda^2\sin 2\beta}, \qquad (39)$$

with the usual W and Z masses

$$m_W^2=\frac{1}{2}g^2v^2 \quad\text{and}\quad m_Z^2=\frac{1}{2}(g^2+g'^2)v^2. \qquad (40)$$

The charged Higgs boson states have mass

$$m_{H^\pm}^2=\frac{2m_s^2}{\sin 2\beta}. \qquad (41)$$

The singlet state N has both scalar and pseudo-scalar states that are degenerate (within our simplifying assumption of $A=C=0$),

$$m_{N_1}^2 = m_{N_2}^2 = \lambda^2 v^2 + m_0^2, \tag{42}$$

while the pseudo-scalar from the Higgs doublets has mass

$$m_{A^0}^2 = \lambda^2 v^2 + m_{H^\pm}^2. \tag{43}$$

The neutral scalars from the doublets have a mass matrix

$$\begin{pmatrix} (\lambda^2 v^2 + m_{H^\pm}^2)\cos^2\beta & (\lambda^2 v^2 - m_{H^\pm}^2)\sin\beta\cos\beta \\ (\lambda^2 v^2 - m_{H^\pm}^2)\sin\beta\cos\beta & (\lambda^2 v^2 + m_{H^\pm}^2)\sin^2\beta \end{pmatrix}, \tag{44}$$

where the upper (lower) components correspond to the h_u^0 (h_d^0) defined by the expansion $H_{u,d}^0 = \langle H_{u,d}^0 \rangle + h_{u,d}^0/\sqrt{2}$. This mass matrix leads to the eigenvalues

$$m_{H^0}^2 = \frac{\lambda^2 v^2 + m_{H^\pm}^2 + X}{2}, \tag{45}$$

$$m_{h^0}^2 = \frac{\lambda^2 v^2 + m_{H^\pm}^2 - X}{2}, \tag{46}$$

where

$$X = \sqrt{(\lambda^2 v^2 + m_{H^\pm}^2)^2 - 4\lambda^2 v^2 m_{H^\pm}^2 \sin^2 2\beta}. \tag{47}$$

They are given in terms of the gauge eigenstates by

$$\begin{pmatrix} h^0 \\ H^0 \end{pmatrix} = \begin{pmatrix} \cos\alpha & -\sin\alpha \\ \sin\alpha & \cos\alpha \end{pmatrix} \begin{pmatrix} h_u^0 \\ h_d^0 \end{pmatrix}, \tag{48}$$

where α is the mixing angle which in our model is

$$\tan\alpha = \frac{(\lambda^2 v^2 + m_{H^\pm}^2)\cos 2\beta + X}{(\lambda^2 v^2 - m_{H^\pm}^2)\sin 2\beta}. \tag{49}$$

The above expressions will receive corrections from the D terms at the level of $m_Z^2/m_{\rm SUSY}^2$ and $m_Z^2/\lambda^2 v^2$.

Here we highlight some of the interesting general features of this spectrum. The model contains singlet scalar and pseudo-scalar states N_1 and N_2 that are not present in the MSSM. In addition, the pseudo-scalar Higgs boson A^0 is always heavier than the charged Higgs boson, which is the opposite of the MSSM. When studying the other states, it is useful to consider two limiting cases, $\lambda v \ll m_{H^\pm}$ and $\lambda v \gg m_{H^\pm}$. In the first case, $\tan\alpha \to -\cot\beta$, and hence the lighter eigenstate ''aligns'' with the vacuum. In other words, the lighter neutral Higgs boson h^0 is standard model–like, while the heavier state H^0 does not have couplings to ZZ or W^+W^-. It forms a custodial $SU(2)$ triplet with the charged Higgs bosons, (H^+, H^0, H^-). In the second case, $\tan\alpha \to \cot\beta$. If $\tan\beta = 1$ they are still aligned, and the *heavier* neutral Higgs boson H^0 is standard model–like, and the custodial $SU(2)$ triplet is (H^+, h^0, H^-). For a general $\tan\beta$, however, the eigenstates and the vacuum are not aligned, and both neutral states contain the standard model–like Higgs boson state. In particular, the mixture is maximal if $\beta = 3\pi/8$ ($\tan\beta = 2.414$).

VII. PHENOMENOLOGY

Since the theory is supersymmetric, we expect superpartners just as in the ordinary MSSM. However, there are important distinctions between our fat Higgs model and the MSSM, especially in the Higgs spectrum. This is the main issue we discuss in this section.

A. Spectrum

To study the phenomenology of the minimal supersymmetric fat Higgs model, we pick three points in the parameter space:

	λ	$\tan\beta$	m_s (GeV)
I	3	2	400
II	2	2	200
III	2	1	200.

(50)

m_0 is chosen to be the same as m_s for simplicity; changing m_0 merely brings the mass of the $N_{1,2}$ states up and down independent of the rest of the spectrum (within our simplifying assumption $A=C=0$) (see Fig. 2).

Spectrum I corresponds to the case $m_{H^\pm} > \lambda v$ where the lightest neutral Higgs boson is standard model–like, while the heavier neutral Higgs boson H^0 forms a triplet with the charged Higgs bosons $H^\pm$ under the approximate custodial $SU(2)$ symmetry. The pseudo-scalar Higgs boson A^0 is heavier than both of them. It resembles the spectrum in the MSSM when A^0 is heavy, but the relative ordering of the heavy Higgs boson states is quite different.

Spectrum II has a smaller supersymmetry breaking scale, and both h^0 and H^0 have significant standard model-like Higgs boson content, approximately 75% and 25%, respectively. Such a large mixing is unusual in the MSSM when the masses are this different.

Spectrum III is the most unconventional of all. Because of the exact custodial $SU(2)$, the triplet h^0 and $H^\pm$ are degenerate, and they do not contain the standard model-like Higgs boson component. On the other hand, the heavier neutral Higgs boson H^0 is standard model-like. The pseudo-scalar Higgs boson is even heavier.

B. Electroweak constraints

It is well known that as the SM-like Higgs mass is raised above about 250 GeV the SM without new physics is increasingly disfavored by electroweak precision data. In our fat Higgs model, however, there are several contributions to electroweak observables that are around the same size as the one from a heavier SM-like Higgs boson. We have calculated the contribution of the Higgs boson states to the electroweak

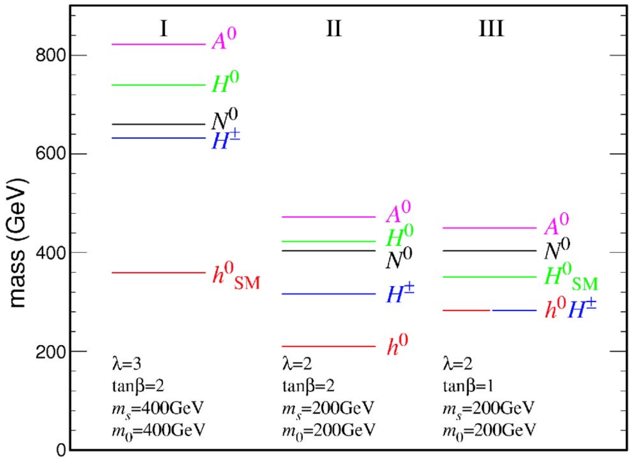

FIG. 2. Sample Higgs spectra in our fat Higgs model. In spectrum I the SM-like Higgs boson is dominantly h^0 (89%), whereas in spectrum III the SM-like Higgs boson is purely H^0.

parameters S, T [27]. The analytical results are the same as the MSSM, and we present formulas in Appendix A for completeness.

We find that the model is consistent with the experimentally allowed region in the S-T plane, described in Appendix B, with no fine-tuning of model parameters. As an example of this, we present the S and T contributions for three trajectories in parameter space in Fig. 3.

In two trajectories, the coupling $\lambda(v_0)$ is varied between 2 and 3 for $\tan\beta=2$ and for two different SUSY breaking scales. The mass of the lightest component of the SM Higgs boson is shown at the end points of each trajectory. Spectrum I of Fig. 2 is at the top of the solid trajectory and spectrum II is at the bottom of the dashed one. Note that these two spectra are in excellent agreement with electroweak constraints despite the heavy Higgs boson masses of 360 and 210 GeV. Furthermore, we find that the constraints for S and T are easily satisfied for a significant range of the model parameters, as the two trajectories demonstrate.

In the third (dotted) trajectory, $\lambda(v_0)$ is varied between 2 and 3 for $\tan\beta=1$. Spectrum III in Fig. 2 is at the top of this line. This trajectory lies mostly within the 99% C.L. contour despite the unconventional Higgs spectrum. However, it is well known that a top-bottom squark splitting $(m^2_{\tilde{t}_L}=m^2_{\tilde{b}_L}+m_t^2)$ may contribute significantly to T (see Fig. 4). This contribution may easily be 0.1–0.5 (e.g. [28]) which would bring these points back into the 68% C.L. ellipse. The size of this contribution depends on the masses in the top-bottom squark system which can generically be different from the SUSY breaking masses in the Higgs sector.

Nevertheless we stress that even with a negligible contribution to T from the top-bottom squark sector, the Higgs boson mass can be hundreds of GeV, as shown by trajectories I and II, and yet still stay well within the precision electroweak contours.

C. $b\to s\gamma$ constraint

Another constraint on our model comes from $b\to s\gamma$ transitions mediated by charged Higgs bosons. Considering only the charged Higgs boson/top quark contribution, the constraint on the charged Higgs boson mass is $m_{H^\pm}>350$ GeV at 99% C.L. [29]. There are two ways this constraint could be satisfied. The first is to simply raise the charged Higgs boson mass above the bound. The second is the well-known possibility that the charged Higgs boson contribution cancels against the chargino–top-squark contribution, and therefore allows a lighter charged Higgs boson [30]. This of course depends on the specific model and parameters for supersymmetry breaking.

D. Search strategies

The neutral Higgs scalars h^0 and H^0 each may have a significant component of the standard model–like Higgs boson and can thus be discovered by the standard search methods, in particular the "gold-plated" signal $h_{SM}\to ZZ\to 4\ell$ at the CERN Large Hadron Collider (LHC). However, the decay modes $H^0\to h^0h^0, H^+H^-$ may also be open, and their partial decay widths are all comparable and proportional to $\lambda^2 m_{H^0}$.

In the strict custodial $SU(2)$ limit, when $\tan\beta=1$, the triplet Higgs bosons are produced only in pairs. In particular, there is no production process of the neutral state by itself, such as $e^+e^-\to Zh$ or $u\bar{d}\to W^+h$, when h does not contain the standard model-like Higgs boson component. Nevertheless, they do have the top Yukawa coupling and thus can be produced from gluon fusion at the LHC. Their decays depend very sensitively on the superpartner spectrum and the Higgs spectrum.

In order to positively establish that our model correctly describes the Higgs sector, numerous other measurements will be needed: the complete mass spectrum, branching fractions, and Higgs boson self-couplings. For this purpose, an

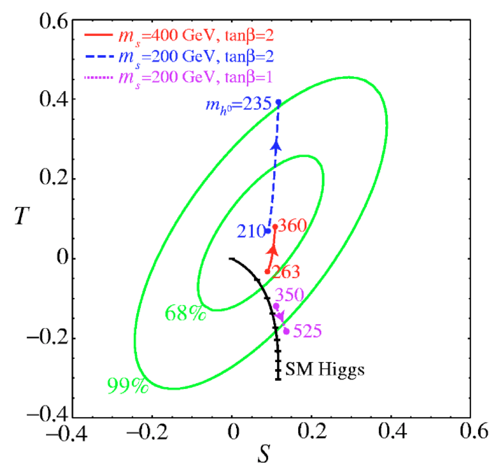

FIG. 3. Constraints on S and T parameters from precision electroweak data at 68% and 99% confidence levels. The plot assumes $U=0$. Contributions of the fat Higgs model to S and T are shown along three trajectories where λ is varied from 2 to 3 in the direction of the arrow. The end points are labeled with the mass (in GeV) of the lighter component of the SM-like Higgs boson. For comparison, the black line shows the contributions to S and T for the standard model with various Higgs boson masses between 100 GeV and 1 TeV in increments of 100 GeV.

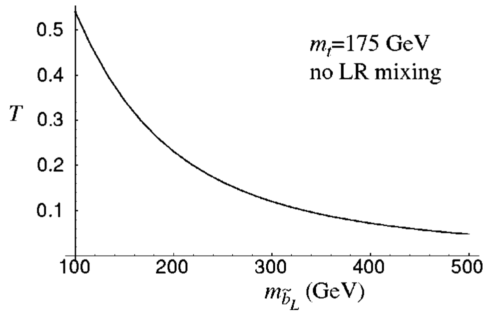

FIG. 4. Contribution of the top-bottom squark sector to the T parameter.

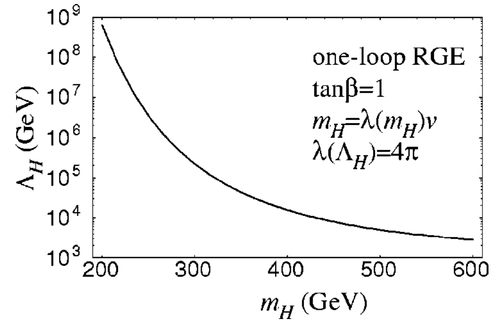

FIG. 5. The compositeness scale is shown as a function of the SM-like Higgs boson mass. For this plot we took $\tan\beta=1$, so the SM-like Higgs boson is purely the heavier state H^0 with mass $m_{H^0}=\lambda v$ (see the discussion at the end of Sec. VI). The scale Λ_H was determined by evolving $\lambda(m_H)$ up to $\lambda(\Lambda_H)=4\pi$ using one-loop renormalization group evolution.

e^+e^- linear collider would be a great asset.

Once the Higgs boson mass is measured, we know its quartic coupling λ. Because of the renormalization group evolution, a lower compositeness scale corresponds to a heavier Higgs boson mass. This is shown in Fig. 5. The limit of the low compositeness scale is of course of special interest, where we may have direct access to the composite dynamics of the Higgs boson. Because this limit has its own special issues, we will defer the discussion of this case, the "fattest Higgs model," to a separate paper.[4]

E. Cosmology

The NMSSM is known to have a cosmological problem due to the spontaneous breaking of its Z_3 symmetry which produces domain walls (see [31] for a recent discussion on this issue). Interestingly, our fat Higgs model does not contain such a symmetry and is free from the domain wall problem. The Z_3 symmetry used to forbid the linear terms in S and S' acts on $T^{1,2}$ with charge $+1$ and $T^{3,4}$ with charge -1, and hence all Higgs fields N, $H_{u,d}$, quarks, and leptons are neutral under this Z_3. On the other hand, R parity can be imposed consistently and thus the lightest supersymmetric particle is a candidate for cold dark matter.

Finally, note that the charge assignments given in Table I for $T^{5,6}$ and the P's and Q's lead to fractionally charged spectators with electric charge $\pm 1/2$. The lightest stable one does not decay, and this could lead to problems in early universe cosmology. There are several options. The first is to simply assume that these particles are not produced after reheating by restricting the reheat temperature to be much lower than their mass. A second possibility is to change the charge assignments of $T^{5,6}$ so that they carry $\pm 1/2$ hypercharge, and therefore all the low energy composites have integral charge. We will see below, however, that leaving the charge assignment as given in Table I allows the simplest interpretation of gauge coupling unification in the model.

VIII. UNIFICATION

In this section we complete the discussion of our fat Higgs model by showing that gauge coupling unification can be easily preserved despite the composite nature of the Higgs fields and the strong coupling. The effective theory below the compositeness scale (and below M_f and $m_{\rm spect}$) has the same matter content as the NMSSM; thus the gauge couplings run exactly like the MSSM gauge couplings until the compositeness scale is reached. That the couplings can unify above the compositeness scale is nontrivial, and to show this we will step through each contribution to the beta functions in the high energy theory (well above the compositeness scale).

The one-loop beta functions for the SM gauge couplings are[5]

$$\frac{d}{dt}g_a=(b_a^{\rm MSSM}+\Delta b_a)g_a^3, \tag{51}$$

where $b_a^{\rm MSSM}$ are the MSSM contributions and the Δb_a characterize the differences between our model and the MSSM. Above the compositeness scale our model has no fundamental Higgs fields. However, the $SU(2)_H$ doublets $T^1,\ldots,T^4$ give exactly the same contribution to the beta functions as do the two Higgs doublets of the MSSM. Thus the selection of $SU(2)_H$ as the strong gauge group has the interesting side effect that the fundamental and composite states give precisely the same contribution to the SM gauge beta functions. Since $T^5,\ldots,T^8$ are neutral under the SM, they do not contribute to the running of the SM gauge couplings.

Our model has two new sectors that contribute to Δb_a: (1) the P and Q fields that marry off spectator composite fields, and (2) the extra doublets $\varphi_{u,d},\overline{\varphi}_{u,d}$ needed to generate fermion masses and mixings.

The first contribution consists of the P and Q fields, yielding

$$\Delta b_1=\frac{3}{5},\quad \Delta b_2=1,\quad \Delta b_3=0, \tag{52}$$

[4]The limit $m\to m'$ is identical to technicolorful supersymmetry [9] except for the presence of the P and Q fields. Our fat Higgs model is hence an "analytic continuation" of technicolorful supersymmetry.

[5]We use the $SU(5)$ GUT normalization $b_1=(3/5)b_Y$.

which corresponds to two $SU(2)_L$ doublets and four fields with hypercharge $\pm 1/2$.

The second contribution, from the extra doublets $\varphi_{u,d},\overline{\varphi}_{u,d}$ needed for fermion masses, is

$$\Delta b_1 = \frac{6}{5}, \quad \Delta b_2 = 2, \quad \Delta b_3 = 0, \tag{53}$$

which corresponds to four $SU(2)_L$ doublets with hypercharge $\pm 1/2$.

The total of Eqs. (52) and (53) is

$$\sum \Delta b_1 = \frac{9}{5}, \ \sum \Delta b_2 = 3, \ \sum \Delta b_3 = 0. \tag{54}$$

Coupling unification requires us to add additional matter at the $\Lambda_H \sim m' \sim M_f$ scale. For instance, we can add three vector-like pairs of chiral multiplets $D_i(\overline{D}_i)$ $(i=1,2,3)$, with the quantum numbers of the right-handed down quarks, i.e. triplets under $SU(3)$ with $U(1)_Y$ quantum numbers $\pm 1/3$. Then

$$\sum \Delta b_1 = \sum \Delta b_2 = \sum \Delta b_3 = 3. \tag{55}$$

This is the same result obtained for a gauge mediation model with three sets of $\mathbf{5}+\overline{\mathbf{5}}$ messengers.

Like gauge mediation, these extra fields have the appearance of "completing" the would-be incomplete $SU(5)$ matter representations. For example, a gauge mediation model with only messenger quark doublets $Q_m + \overline{Q}_m$ is sufficient to communicate supersymmetry breaking, but as these fields do not form a complete $SU(5)$ representation, additional messenger fields filling up the $\mathbf{10}_m$ and $\overline{\mathbf{10}}_m$ must be added to preserve gauge coupling unification. Unlike gauge mediation, however, the extra color triplets $D(\overline{D})$ cannot be in the same GUT representation as $\varphi_{u,d},\overline{\varphi}_{u,d}$, otherwise dimension-6 triplet-induced proton decay will be too fast. In any case, we have shown that adding three pairs of color triplets to our model does not affect the dynamics and yet provides an existence proof that gauge coupling unification can work just as well as in the MSSM.

It is also important to emphasize that we do not expect large threshold corrections from passing through the strong coupling or superconformal sector. Due to holomorphy, the low-energy gauge couplings are determined only by the bare mass of the heavy particles that are integrated out [32]. This can also be seen by noting that in supersymmetric theories both the exact Novikov-Shifman-Vainshtein-Zakharov beta function [33,34] and the decoupling mass depend on the wave-function renormalization factor, which drops out from the final result. Therefore, gauge coupling unification is unaffected even in the presence of strong $SU(2)_H$ dynamics in which the standard model gauge groups $SU(3)_c \times SU(2)_L \times U(1)_Y$ are perturbatively coupled. The dominant effect is therefore the threshold correction resulting from potential differences between the mass of the color triplets, the mass $m_{\rm spect}$ of the spectators, and the mass M_f of the extra doublets φ. Suppose that the same flavor symmetry that ensures the T^7T^8 doublets acquire the mass m' could also be used to determine the color triplet masses. In this case the threshold corrections are no larger than $\log m'/M_f$ or $\log m'/m_{\rm spect}$, of the same order of magnitude as the MSSM or GUT threshold corrections, which is much smaller than the leading $\log M_{\rm unif}/M_Z$ in the MSSM.

We have shown that gauge coupling unification can be preserved with a small number of additional matter fields, but it is obvious that we cannot embed the matter content into a single four-dimensional GUT group. Unification of the gauge couplings therefore could be due to string unification or orbifold GUT unification in five [35] (or four [36]) dimensions, where the matter content does not need to fall into a GUT representation [37].

IX. DISCUSSION

We have constructed a supersymmetric composite Higgs boson theory that solves the supersymmetric little hierarchy problem. Electroweak symmetry is broken dynamically through a new gauge interaction that gets strong at an intermediate scale. The composite Higgs fields have a dynamically generated superpotential that has a form similar to the NMSSM, and hence solves the μ problem, but with no restriction on the coupling λ. This allows the tree-level Higgs boson mass to be much higher, 200–450 GeV, solving the supersymmetric little hierarchy problem. The usual lore about upper bounds on the lightest Higgs boson mass in supersymmetric theories is therefore obviously violated. With hindsight we see that requiring perturbativity of the Higgs sector was simply too restrictive. To the best of our knowledge, the fat Higgs model provides the first explicit example where the Higgs sector is composite and yet the dynamics are fully calculable and UV complete.

There are several interesting future avenues of research. We used the Giudice-Masiero mechanism to determine certain mass scales and therefore supergravity mediation was implicit. For generic choices of the supergravity-mediated contributions we have the regular supersymmetric flavor-changing neutral current problem. One solution is a flavor symmetry, e.g. $U(3)^5$ in [38]. Another possibility is to implement one of several flavor-blind supersymmetry breaking mechanisms such as gauge mediation [39], anomaly mediation [40] [supplemented by $U(1)$ D terms to make it viable with UV insensitivity [41]] or its 4D realization [42], or gaugino mediation in five [43] or four [44] dimensions. It remains to be seen whether these mediation mechanisms achieve an acceptable mass spectrum and electroweak symmetry breaking. These methods of supersymmetry breaking would also require a different mechanism to naturally determine the scales. It would also be interesting to explore unification further in this model, such as whether $SU(2)_H$ can be unified with the other SM gauge groups.

Finally, we have shown that the Higgs boson mass spectrum is quite unusual. It is important to study specifically how our fat Higgs model can be distinguished from more conventional supersymmetric models at future collider experiments. Clearly more work is needed. We cannot overemphasize the importance of next generation experiments being able to analyze their data with as few theoretical assumptions as possible.

ACKNOWLEDGMENT

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APPENDIX A: HIGGS SECTOR CONTRIBUTION TO *S* AND *T*

Here we provide expressions for the perturbative contribution of the Higgs sector to S and T. The Higgs sector consists of mass eigenstates $H^{\pm}$, H^0, h^0, and A^0 which fit into an $SU(2)_L$ doublet,

$$\begin{pmatrix} H^+ \\ \frac{1}{\sqrt{2}}(\tilde{H}^0+iA^0) \end{pmatrix}, \tag{A1}$$

where $\tilde{H}^0=\cos(\beta-\alpha)h^0-\sin(\beta-\alpha)H^0$, and a standard model–like neutral scalar $\tilde{h}^0$, which is orthogonal to $\tilde{H}^0$. The Higgsinos get a vector-like mass of order $m_{\rm SUSY}$, so their contribution to S and T is negligible. $\tilde{h}^0$ will contribute much like a heavy standard model Higgs boson. However, since it is not a pure mass eigenstate, the exact contribution is fairly complicated. We approximate this contribution by taking the weighted sum of the two-loop contributions extracted from the electroweak observables discussed in Appendix B.

Below we summarize the one-loop contribution of the scalar Higgs doublet. Note that the various components of the doublet all have different masses, and the neutral scalar $\tilde{H}^0$ is itself a linear combination of two mass eigenstates. Taking this mixing into account, we find

$$\Delta S=\sin^2(\beta-\alpha)F(m_{H^{\pm}},m_{H^0},m_{A^0}) + \cos^2(\beta-\alpha)F(m_{H^{\pm}},m_{h^0},m_{A^0}), \tag{A2}$$

where the function F is defined by

$$F(m_1,m_2,m_3)=\frac{1}{2\pi}\int_0^1 dx x(1-x)\log\frac{(1-x)m_2^2+xm_3^2}{m_1^2}. \tag{A3}$$

Similarly for T, we find

$$\Delta T=\frac{1}{16\pi m_W^2 s_W^2}[\sin^2(\beta-\alpha)G(m_{H^{\pm}},m_{H^0},m_{A^0}) + \cos^2(\beta-\alpha)G(m_{H^{\pm}},m_{h^0},m_{A^0})], \tag{A4}$$

where the function G is defined as

$$\begin{aligned} G(m_1,m_2,m_3) &= m_2^2 I(m_3,m_2,m_1,m_2)+m_3^2 I(m_2,m_3,m_1,m_3) \\ &\quad -m_1^2 I(m_2,m_1,m_1,m_1)-m_1^2 I(m_3,m_1,m_1,m_1) \end{aligned} \tag{A5}$$

in terms of the integral I,

$$I(m_1,m_2,m_3,m_4)=2\int_0^1 dx x \log\frac{(1-x)m_1^2+xm_2^2}{(1-x)m_3^2+xm_4^2}. \tag{A6}$$

APPENDIX B: *S*-*T* CONTOURS

The experimental constraints on the S-T plane can be easily computed approximately using the following method. We follow the path of Marciano [45] and Perelstein-Peskin-Pierce [46] to focus on only three observables, m_W, Γ_l, and $s_*^2\approx\sin^2\theta_{\rm eff}^{\rm lept}$ from the asymmetries as they are the most accurately measured and sensitive observables to the oblique corrections.

Expressions for these observables, including their approximate m_t, α, and m_H dependence, have been computed by Degrassi and Gambino [47]. We add to those expressions the dependence on S and T as found in Appendix B of [27]. The LEP Electroweak Working Group recommends $\Delta\alpha_{\rm had}^{(5)}(m_Z^2)=0.02761(36)$, including the BES data as discussed in Sec. 16.3 of [48], which implies $\alpha^{-1}(m_Z)=128.945\pm 0.049$. Expanding to linear order in Δm_t and $\Delta\alpha^{-1}$ about $m_t=174.3$ GeV and $\alpha^{-1}=128.945$ leads to the expressions

$$\begin{aligned} m_W &= 80.380+0.13\Delta\alpha^{-1}+0.0061\Delta m_t-0.29S+0.44T \\ &\quad +0.34U-0.058l_h-0.008l_h^2, \\ s_*^2 &= 0.23140-0.0026\Delta\alpha^{-1}-0.000032\Delta m_t+0.0036S \\ &\quad -0.0025T+0.00052l_h, \\ \Gamma_l &= 84.011+0.12\Delta\alpha^{-1}+0.009\Delta m_t-0.19S+0.78T \\ &\quad -0.054l_h-0.021l_h^2, \end{aligned} \tag{B1}$$

where $l_H=\log(m_H/100\ \text{GeV})$.

For the experimental values, we use [46]

$$m_W=(80.425\pm 0.034)\ \text{GeV}, \tag{B2}$$

$$s_*^2=0.23150\pm 0.00016, \tag{B3}$$

$$\Gamma_l=(83.984\pm 0.086)\ \ \text{MeV}. \tag{B4}$$

Then χ^2 is defined as

$$\chi^2=\frac{(m_W-80.425)^2}{0.034^2}+\frac{(s^2-0.23150)^2}{0.00016^2}+\frac{(\Gamma_l-83.984)^2}{0.086^2}+\frac{(\Delta\alpha^{-1})^2}{0.049^2}+\frac{(\Delta m_t)^2}{5.1^2}, \tag{B5}$$

which is first minimized with respect to $\Delta\alpha^{-1}$ and Δm_t for each (S,T). This expression for χ^2 yields contours that agree very well with those by the Particle Data Group [49].

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