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A fat Higgs with a fat top

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Abstract: A new variant of the supersymmetric fat Higgs model is presented in which the MSSM Higgses as well as the top quark are composite. The underlying theory is an s-confining SU(3) gauge theory with the MSSM gauge groups realized as gauged sub-groups of the chiral flavor symmetries. This motivates the large Yukawas necessary for the large top mass and SM-like Higgs of mass $\gg M_Z$ in a natural way as the residual of the strong dynamics responsible for the composites. This removes fine-tuning associated with these couplings present in the original fat Higgs and "new fat Higgs" models, respectively.

Keywords: Supersymmetric Effective Theories, Supersymmetric Standard Model, Beyond Standard Model.

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1. Introduction

Supersymmetry (SUSY) is the most cherished and best studied vision of physics beyond the Standard Model (SM). SUSY tames the quadratic divergences that destabilize the electroweak (EW) scale, and results in a host of new particles which should be discovered in the near future if the SUSY vision of particle physics should prove correct.

However, LEP-II has left the minimal supersymmetric standard model (MSSM) in an interesting situation [1]. The minimal model predicts a light Higgs whose tree-level mass is at most M_Z, in contradiction with the LEP-II limit of $M_h^{(SM)} \geq 115\,\mathrm{GeV}$. In order to survive the LEP limit, one must either invoke very large radiative corrections from the top sector [2], CP violation chosen in a very particular way [3], or abandon the minimal model in favor of more ingredients [4−10]. The invocation of large radiative corrections is particularly troublesome, because this tends to introduce unacceptably large corrections to the EW scale, recreating a "little hierarchy problem". While there is some uncertainty in the estimates for the lightest CP even Higgs mass originating in the uncertainty in measured top mass, it appears that the MSSM requires fine-tuning at the level of a few per cent if it is to be consistent with LEP data, and is uncomfortably fine-tuned. This is the "Supersymmetric Little Hierarchy Problem".

The fat Higgs (FH) [7] is a particular, interesting solution to this dilemma. It proposes an alternative to the standard MSSM picture of electroweak symmetry breaking (EWSB) and results in a heavier "light" CP-even Higgs than can be realized in that standard scenario, thus naturally evading the LEP-II bounds. It originates from an s-confining theory, in which a number of fundamental preons charged under a strong SU(2) form Higgs bosons as composites. A variation [8] has a composite singlet from an s-confining SU(4) theory, but the EWSB Higgses are fundamentals. Both theories have interesting distinctive SUSY Higgs phenomenology [6, 7], largely due to the fact that the Higgs quartic interaction may be much larger than is suggested by perturbative unification [11].

Both of these FH theories are challenged in producing large Yukawa interactions. The original FH must generate fermion masses through Yukawa interactions which couple the composite H and $\overline{H}$ to the fundamental quarks and leptons. At the level of the preons, this is a non-renormalizable super-potential coupling, which the original FH generates from renormalizable interactions by integrating out a pair of Higgs-like fields uncharged under the strong SU(2) (see figure 1). The resulting Yukawas thus depend on fundamental parameters as,

$$y_{\text{eff}} \sim \frac{yy'}{4\pi}\frac{\Lambda}{M_H} \tag{1.1}$$

in which y, y' are Yukawas between the preons and/or fundamental fermion superfields (at the compositeness scale Λ), 4π is the naive dimensional analysis (NDA) counting [12] for the coupling of a composite to fundamental fields, Λ is the scale of s-confinement of the strong SU(2) and M_H is the (supersymmetric) mass of the Higgs-like fields. For the light fermions, this is not problematic. Small fermion masses are easily realized. For the top quark, producing a Yukawa coupling of order one requires tuning the scales Λ and M_H to be close to one another (which is somewhat counter-intuitive since they are in principal unrelated to one another, though it was argued in [7] that the coincidence of scales could arise from a flavor symmetry) and that the underlying y and y' be large at Λ to compensate for the 4π. This last fact is also potentially a source of fine-tuning. The strong SU(2) tries to renormalize y and y' strong at low energies. This is helpful in that it compensates the suppression, but dangerous because a large super-potential coupling may ruin the conformal regime of the theory above Λ.

While it is possible that interesting (and phenomenologically viable) low energy dynamics would emerge in this case, the additional strong y (and/or in generalizations y') couplings potentially disrupt the low energy s-confinement solution, and makes it difficult to draw firm conclusions about the low energy physics. One is thus forced to assume that y and y' become moderately strong, but do not quite reach truly strong coupling before the s-confinement scale. Another way to consider the tension is to note[1] that one must

[1] We are indebted to Kaustubh Agashe for discussions on this point.

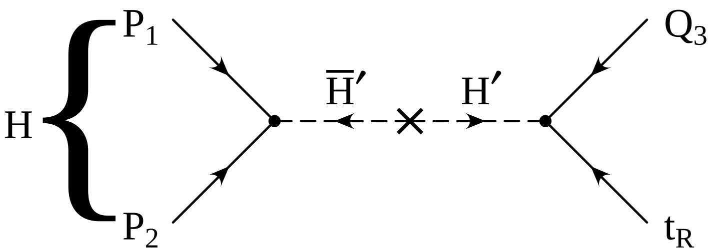

Figure 1: Example graph for how the top Yukawa coupling is generated in the fat Higgs model by integrating out a pair of Higgs-like superfields ($\overline{H}'$, H') to generate a non-renormalizable interaction between preons (P_1 and P_2) bound into a composite Higgs H.

tune the original y and y' to some very particular values in the UV such that they become large enough (but not too large) at Λ. The "new fat Higgs" [8] avoids this issue for the top Yukawa, because in that case the EW Higgses and the quarks are fundamental. Thus, the strong SU(4) does not effectively drive that interaction strong at low energies. However, it recreates the problem for the Higgs quartic itself, because now the quartic links the composite EW singlet S to the fundamental EW Higgses H and $\overline{H}$, and thus feels the same sort of tension when one tries to obtain a large Higgs quartic.

In this article, we explore a new incarnation of the fat Higgs. Our theory is an $SU(3)_s$ SUSY gauge theory which s-confines, producing a composite singlet S and doublets H and $\overline{H}$ as in the original fat Higgs. However, the additional preons are arranged such that they also produce a composite third generation quark doublet (Q_3) and up-type singlet (t_R). The dynamically generated super-potential contains the terms needed for FH-style EWSB, but it also includes the top Yukawa coupling. Since all fields requiring large Yukawa interactions are composite, we have removed the need for strong underlying Yukawa interactions, and thus the danger that the low energy physics could be spoiled by out-of-control non-perturbative couplings. Furthermore, while we will still need to invoke massive fields to generate the Yukawa interactions of the light fermions, there is considerably less need to fine-tune the mass of these "spectator" superfields (M_H) to the s-confinement scale Λ, and/or invoke underlying super-potential couplings which are dangerously large.

In section 2, we present the model and show how it gives rise to all of the required low energy structure of the MSSM. In section 3, we address some of the issues regarding high energy gauge coupling unification. In section 4 we discuss some of the distinctive phenomenology. And in section 5 we conclude.

2. An SU(3) model

Our model has an extended gauge symmetry,

$$SU(3)_s \times SU(3)_c \times SU(2)_W \times U(1)_Y\,. \tag{2.1}$$

$SU(3)_s$ is a "strong" group which will be responsible for generating the MSSM Higgses, a fat-Higgs like singlet, and top from a set of preons, and the remaining gauge groups are as in the MSSM. The particle content charged under $SU(3)_s$ consists of a set of preons

	$SU(3)_s$	$SU(3)_c$	$SU(2)_W$	$U(1)_Y$	$\mathbb{Z}_2$
P_3	$\square$	$\square$	$\mathbf{1}$	0	$+$
P_1	$\square$	$\mathbf{1}$	$\mathbf{1}$	$-2/3$	$-$
$\overline{P}_2$	$\overline{\square}$	$\mathbf{1}$	$\square$	$+1/6$	$-$
$\overline{P}_1$	$\overline{\square}$	$\mathbf{1}$	$\mathbf{1}$	$+2/3$	$+$
$\overline{P}_{\tilde{1}}$	$\overline{\square}$	$\mathbf{1}$	$\mathbf{1}$	$-1/3$	$-$
P'	$\square$	$\mathbf{1}$	$\mathbf{1}$	$+1/3$	$-$
$\overline{P}'$	$\overline{\square}$	$\mathbf{1}$	$\mathbf{1}$	$-1/3$	$-$

Table 1: The $SU(3)_s$-charged Preons. The first set are those participating in the s-confining phase. The second category are integrated out, triggering s-confinement.

	$SU(3)_s$	$SU(3)_c$	$SU(2)_W$	$U(1)_Y$	$\mathbb{Z}_2$
L_i	$\mathbf{1}$	$\mathbf{1}$	$\square$	$-1/2$	$-$
e_i	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	$+1$	$-$
$Q_{1,2}$	$\mathbf{1}$	$\square$	$\square$	$+1/6$	$-$
d_i	$\mathbf{1}$	$\overline{\square}$	$\mathbf{1}$	$+1/3$	$-$
$u_{1,2}$	$\mathbf{1}$	$\overline{\square}$	$\mathbf{1}$	$-2/3$	$-$
$\overline{q}_1$	$\mathbf{1}$	$\overline{\square}$	$\mathbf{1}$	$-2/3$	$+$
$\overline{q}_2$	$\mathbf{1}$	$\overline{\square}$	$\mathbf{1}$	$+1/3$	$-$
H'	$\mathbf{1}$	$\mathbf{1}$	$\square$	$+1/2$	$+$
$\overline{H}'$	$\mathbf{1}$	$\mathbf{1}$	$\square$	$-1/2$	$+$

Table 2: Additional fundamental fields for the SU(3) model. The index $i = 1, 2, 3$ denotes the usual generation number.

listed in table 1. Since the matter is vector-like with respect to $SU(3)_s$, we follow the usual fashion and refer to it as a "SUSY QCD" theory, but this should not be confused with the usual color interaction of the MSSM, $SU(3)_c$. Note that the MSSM gauge groups are gauged sub-groups of the $SU(F) \times SU(F) \times U(1)_B$ chiral symmetries. The set of preons is non-anomalous (in fact, it is vector-like) with respect to $SU(3)_s$, and there are no mixed anomalies between $SU(3)_s$ and the MSSM gauge groups. However, the MSSM gauge symmetries are anomalous with respect to themselves. This is in fact related to the point that the strong sector will eventually give rise to a composite Q_3, t_R, $\overline{H}$, S and H, but not to b_R, L_3, or e_3. Thus, we introduce a set of fundamental fields uncharged under $SU(3)_s$ in table 2. The first and second generation superfields appear as fundamental fields, as in the MSSM. Also indicated are the charges of the fields under a $\mathbb{Z}_2$ "R-parity" which plays the same role to suppress dangerous renormalizable baryon- and lepton-number violating processes as it does in the MSSM. The assignment of preon hypercharges is not completely determined by requiring the correct hypercharges for the composites, and the particular choice we make is based partly on aesthetics (requiring that all exotic colored

		SU(3)$_c$	SU(2)$_W$	U(1)$_Y$	$\mathbb{Z}_2$
$B_1 \leftrightarrow t_R$	$P_3 P_3 P_1$	$\overline{\square}$	$\mathbf{1}$	$-2/3$	$-$
$B_2 \leftrightarrow S$	$P_3 P_3 P_3$	$\mathbf{1}$	$\mathbf{1}$	0	$+$
$\overline{B}_1 \leftrightarrow H$	$\overline{P}_2 \overline{P}_1 \overline{P}_{\tilde{1}}$	$\mathbf{1}$	$\square$	$+1/2$	$+$
$\overline{B}_2 \leftrightarrow \psi$	$\overline{P}_2 \overline{P}_2 \overline{P}_1$	$\mathbf{1}$	$\mathbf{1}$	$+1$	$+$
$\overline{B}_3 \leftrightarrow \chi$	$\overline{P}_2 \overline{P}_2 \overline{P}_{\tilde{1}}$	$\mathbf{1}$	$\mathbf{1}$	0	$-$
$M_1 \leftrightarrow Q_3$	$P_3 \overline{P}_2$	$\square$	$\square$	$+1/6$	$-$
$M_2 \leftrightarrow q_1$	$P_3 \overline{P}_1$	$\square$	$\mathbf{1}$	$+2/3$	$+$
$M_3 \leftrightarrow q_2$	$P_3 \overline{P}_{\tilde{1}}$	$\square$	$\mathbf{1}$	$-1/3$	$-$
$M_4 \leftrightarrow \overline{H}$	$P_1 \overline{P}_2$	$\mathbf{1}$	$\square$	$-1/2$	$+$
$M_5 \leftrightarrow \overline{\chi}$	$P_1 \overline{P}_1$	$\mathbf{1}$	$\mathbf{1}$	0	$-$
$M_6 \leftrightarrow \overline{\psi}$	$P_1 \overline{P}_{\tilde{1}}$	$\mathbf{1}$	$\mathbf{1}$	-1	$+$

Table 3: Composites of the SU(3) model.

particles have charges $\pm 1/3$ or $\pm 2/3$ and all exotic uncolored particles have charges ± 1 or zero), and partly motivated by gauge coupling unification as we shall see below. Many fundamental Yukawa interactions can be formed out of these fields. To preserve readability, we discuss these in groups in the subsections below.

This theory is SUSY SU(3) QCD with 5 flavors, which is inside the conformal window [13]. From any value of the SU(3)$_s$ gauge coupling at very high scales, it flows (assuming, as we will do so, that all of the fundamental Yukawa interactions are not strong enough to disrupt the approximate scale-invariance) at lower scales to the fixed point at,

$$g_*^2 \simeq \frac{4\pi^2}{3}\,. \tag{2.2}$$

We include a super-potential mass for P' (and for the uncolored H'),

$$W_m = M_P \overline{P}' P' + M_H \overline{H}' H'\,. \tag{2.3}$$

Below M_P, the P', $\overline{P}'$ flavor may be integrated out and the theory loses conformality, flowing to an s-confining phase [14]. We denote the confinement scale by Λ, and estimate from the large fixed point coupling g_* that the two scales are approximately equal,

$$\Lambda \simeq M_P\,. \tag{2.4}$$

The scale M_P must be input by hand, and determines the strong coupling scale Λ.

2.1 Composites and dynamical super-potential

Below the confinement scale, the theory can be described by composite SU(3)$_s$-invariant mesons (M) and baryons (B, $\overline{B}$), listed in Table 3. A dynamical super-potential is generated with form,

$$\begin{aligned} W_{\text{dyn}} &= \frac{1}{\Lambda^5}\left\{\overline{B} M B - \det M\right\} \\ &\to \lambda \left\{ H Q_3 t_R + H \overline{H} S + \psi q_2 t_R + \psi \overline{\psi} S + \chi \overline{\chi} S + \chi q_1 t_R - \frac{\lambda}{\Lambda} \det M \right\}, \end{aligned} \tag{2.5}$$

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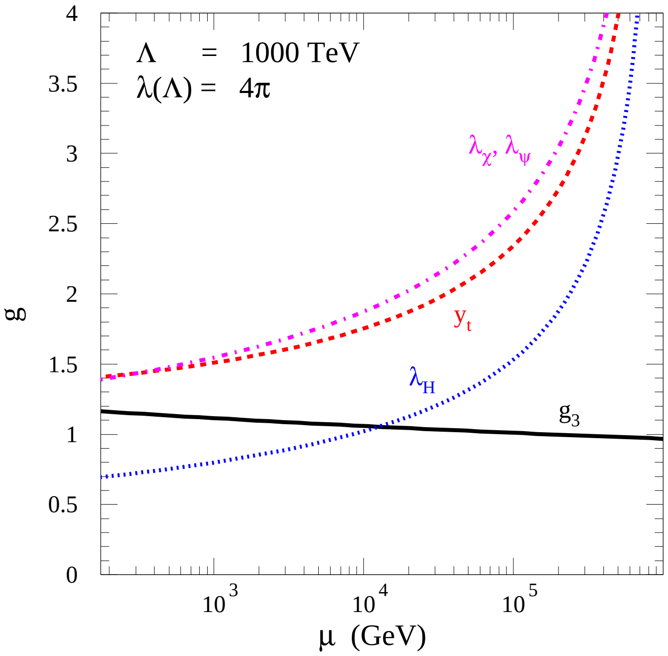

Figure 2: The RGE evolution from $\lambda = 1000\,\text{TeV}$ to v of the strong coupling g_3 (solid curve), top Yukawa interaction y_t (dashed curve), $SH\overline{H}$ interaction λ_H (dotted curve), and $S\psi\overline{\psi}$ and $S\chi\overline{\chi}$ interactions λ_ψ and λ_χ (dot-dashed curve).

where in the second line we rescaled the baryons and mesons to canonically normalized superfields. It will not be very important for our purposes, but we note for completeness that one may express the irrelevant interactions as,

$$\det M = \epsilon_{ij}\epsilon_{\alpha\beta\gamma}\left(\overline{H}^i Q_3^{\alpha j} q_1^\beta q_2^\gamma + \overline{\chi} Q_3^{\alpha i} Q_3^{\beta j} q_2^\gamma + \overline{\psi} Q_3^{\alpha i} Q_3^{\beta j} q_1^\gamma\right), \tag{2.6}$$

suppressed by the confinement scale λ/Λ. We have also provided the naive dimensional analysis (NDA) estimate for the coupling $\lambda \sim 4\pi$ [12]. Thus, this model dynamically generates the fat Higgs sector and super-potential, along with the top Yukawa coupling and some exotic interactions with exotic superfields. Note that the exotics occur in pairs in these interactions, because they arise exclusively from composites which include an odd number of $\overline{P}_1$ and $\overline{P}_{\tilde{1}}$.

We shall see below that q_1 and q_2 receive masses of order Λ. Thus, below Λ the relevant couplings in (2.5) are the top Yukawa y_t, the $SH\overline{H}$ interaction λ_H, the $S\psi\overline{\psi}$ interaction λ_ψ, and the $S\chi\overline{\chi}$ interaction λ_χ. All of these are equal and of order $\lambda \sim 4\pi$ at the scale Λ, but because the $q's$ decouple at that scale, and because of our having gauged subgroups of the chiral symmetries of the SUSY QCD theory, they evolve apart at lower energies.

In order to discuss the top mass and EWSB, these should be evolved down to energy scales of order the electroweak scale v. At one-loop, below Λ, the dominant renormalization effects are from y_t, and $\lambda_{(H,\psi,\chi)}$ themselves, and from the $\mathrm{SU}(3)_c$ coupling g_3. The one loop renormalization group equations (RGEs) are

$$\frac{dg_3}{dt} = -\frac{3}{16\pi^2} g_3^3 \tag{2.7}$$
$$\frac{dy_t}{dt} = \frac{y_t}{16\pi^2}\left[6|y_t|^2 + |\lambda_H|^2 - \frac{16}{3}g_3^3\right] \tag{2.8}$$
$$\frac{d\lambda_H}{dt} = \frac{\lambda_H}{16\pi^2}\left[3|y_t|^2 + 4|\lambda_H|^2 + |\lambda_\psi|^2 + |\lambda_\chi|^2\right] \tag{2.9}$$
$$\frac{d\lambda_\psi}{dt} = \frac{\lambda_\psi}{16\pi^2}\left[2|\lambda_H|^2 + 3|\lambda_\psi|^2 + |\lambda_\chi|^2\right] \tag{2.10}$$
$$\frac{d\lambda_\chi}{dt} = \frac{\lambda_\chi}{16\pi^2}\left[2|\lambda_H|^2 + 3|\lambda_\chi|^2 + |\lambda_\psi|^2\right], \tag{2.11}$$

where t is the renormalization scale $t \equiv \log \mu_R$. Since $\lambda_\psi = \lambda_\chi$ at scale Λ, these coupling strengths will remain equal up to very small effects from the different hypercharges of ψ and χ.

The fact that the top mass has been measured at the Tevatron [15] allows us to approximately fix Λ, up to the choice of $\tan\beta$. As values of $\tan\beta \sim 1$ result in the largest light CP even Higgs masses, we make this choice for which the target y_t is about $\sqrt{2}$. Solving the coupled equations numerically and imposing this requirement fixes $\Lambda \sim 10^4 \times v$ (i.e. $\Lambda \sim 1000\,\mathrm{TeV}$), and predicts that λ_H will be somewhat less than y_t itself. An example is shown in figure 2. Note that there are order one uncertainties in $\lambda(\Lambda)$, which could easily modify our estimate for Λ by an order of magnitude.[2] Irrespectively, the prediction that the Higgs quartic is approximately locked to the top Yukawa interaction is an interesting feature of the model.

2.2 Electroweak symmetry breaking

We include a Yukawa coupling in the fundamental theory,

$$\begin{aligned} W_S &= -y_S \epsilon_{\alpha\beta\gamma} P_3^\alpha P_3^\beta P_3^\gamma \\ &\to -\left(\frac{y_S}{4\pi}\Lambda^2\right) S\,, \end{aligned} \tag{2.12}$$

(where α, β, and γ are $\mathrm{SU}(3)_c$ indices, and the $\mathrm{SU}(3)_s$ indices are similarly contracted anti-symmetrically but not shown for clarity) which becomes a tadpole for S below Λ. Combined with W_{dyn}, this results in Higgs super-potential,

$$W_H = \lambda_H S\left(H\overline{H} - v_0^2\right) + \lambda_\psi S\psi\overline{\psi} + \lambda_\chi S\chi\overline{\chi}\,, \tag{2.13}$$

where v_0^2 has NDA estimate (at scale Λ),

$$v_0^2 \sim \frac{y_S}{\lambda\,(4\pi)}\Lambda^2 \sim \frac{y_S}{(4\pi)^2}\Lambda^2 \tag{2.14}$$

thus indicating that v_0 is naturally at least an order of magnitude below Λ, and will be smaller if y_S takes a sufficiently small value (as we will assume it does in order to

[2] There are also order one uncertainties in the RGE evolution from higher orders close to scale Λ, where the couplings are strong, as well.

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appropriately generate the EW scale). Aside from the presence of the additional superfields ψ, $\overline{\psi}$, χ, $\overline{\chi}$, this is the super-potential of the fat Higgs, leading to a electroweak symmetry-breaking even in the supersymmetric limit.

The scalar Higgs potential consists of the contribution from the dynamical super-potential above, the MSSM D-terms, and the corrections from soft SUSY breaking. There is also an effective μ term induced by integrating out H' and $\overline{H}'$ as described below in section 2.3.4. Altogether, this leads to a scalar potential,

$$\begin{aligned} V_H = & |\lambda_H H\overline{H} + \lambda_\psi \psi\overline{\psi} + \lambda_\chi \chi\overline{\chi} - v_0^2|^2 + \lambda_H^2 \left(|SH|^2 + |S\overline{H}|^2\right) + \\ & + \lambda_\psi^2 \left(|S\psi|^2 + |S\overline{\psi}|^2\right) + \lambda_\chi^2 \left(|S\chi|^2 + |S\overline{\chi}|^2\right) + \\ & + \frac{g_2^2}{8}\left(H^\dagger \vec{\tau} H + \overline{H}^\dagger \vec{\tau}\overline{H}\right)^2 + \frac{g_1^2}{2}\left(\frac{1}{2}|H|^2 - \frac{1}{2}|\overline{H}|^2 + |\psi|^2 - |\overline{\psi}|^2\right)^2 + \\ & + \left(m_H^2 + |\mu|^2\right)|H|^2 + \left(m_{\overline{H}}^2 + |\mu|^2\right)|\overline{H}|^2 + m_S^2|S|^2 + \\ & + m_\psi^2|\psi|^2 + m_{\overline{\psi}}^2|\overline{\psi}|^2 + m_\chi^2|\chi|^2 + m_{\overline{\chi}}^2|\overline{\chi}|^2 + \\ & + \left\{A_S\left(\lambda_H S H\overline{H} + \lambda_\psi S\psi\overline{\psi} + \lambda_\chi S\chi\overline{\chi}\right) - T_S v_0^2 S + h.c.\right\}, \end{aligned} \tag{2.15}$$

where $g_{1,2}$ are the MSSM U(1)/ SU(2) gauge couplings, and the m's, A_S, and T_S are soft SUSY breaking parameters. We have assumed that the A terms are locked together by the underlying chiral symmetries of the SUSY QCD theory, and in the same spirit ignored other potential SUSY breaking terms such as $B\mu$-like terms involving $H\overline{H}$, $\psi\overline{\psi}$, and $\chi\overline{\chi}$. Of course, we expect that the equality of the A terms is only approximate, as the RGEs will split them apart just as it does the λ interactions, but we continue to neglect such splittings to simplify the discussion.

In general, the minimization conditions are quite complicated, but we sketch a solution below. To simplify matters, we begin by considering $m_H = m_{\overline{H}} = m_S \equiv m$, $m_\psi = m_{\overline{\psi}} = m_\chi = m_{\overline{\chi}} \equiv M$, $A_H = A_\psi = A_\chi = T_S = 0$, and ignore the MSSM D-terms. We will consider deviations from these assumptions below. Under these conditions, the potential is symmetric under $H \leftrightarrow \overline{H}$ and $\psi \leftrightarrow \overline{\psi} \leftrightarrow \chi \leftrightarrow \overline{\chi}$. The SM-like Higgs is $h = (H^0 + \overline{H}^0)/\sqrt{2}$, and we denote the common vacuum expectation value (VEV) of ψ, $\overline{\psi}$, χ, and $\overline{\chi}$ as $\phi/\sqrt{2}$. The scalar potential becomes

$$\begin{aligned} & \left(\frac{\lambda_H^2}{4}h^4 + \lambda_\psi\phi^4 + \lambda_H\lambda_\psi h^2\phi^2 + 2\lambda_\psi^2|S|^2\phi^2\right) + m^2|S|^2 + \\ & + \left(m^2 + |\mu|^2 - \lambda_H^2 v_0^2\right)h^2 + \left(M^2 - \lambda_\psi^2 v_0^2\right)\phi^2 \end{aligned} \tag{2.16}$$

and the vacuum crucially depends on the signs of the quantities $(m^2 + |\mu|^2 - \lambda_H^2 v_0^2)$ and $(M^2 - \lambda_\psi^2 v_0^2)$. Under the relatively mild requirement that the soft masses respect,

$$\left(m^2 + |\mu|^2 - \lambda_H^2 v_0^2\right) < 0 \tag{2.17}$$

$$\left(M^2 - \lambda_\psi^2 v_0^2\right) > 0 \tag{2.18}$$

we arrive at the solution $\langle H\rangle = \langle\overline{H}\rangle = \sqrt{v_0^2 - (m^2 + |\mu|^2)/\lambda_H^2}$, $\langle S\rangle = \langle\psi\rangle = \langle\overline{\psi}\rangle = \langle\chi\rangle = \langle\overline{\chi}\rangle = 0$, leading to viable[3] EWSB. Including the D terms and relaxing the universality

[3]Note that a VEV for ψ or $\overline{\psi}$ would lead to large (tree level) corrections to $\Delta\rho$.

among the soft masses will not disrupt this general feature, provided m_ψ, $m_{\overline{\psi}}$, m_χ, and $m_{\overline{\chi}}$ continue to individually satisfy eq. (2.18), though it will modify the expressions for the VEVs and cause $\tan\beta \equiv \langle H\rangle / \langle \overline{H}\rangle$ to deviate from unity.

We also consider non-zero values for A_S and T_S. Both of these terms, combined with the EWSB VEVs for H and $\overline{H}$, generate tadpoles for S which will generically result in it acquiring a VEV of order the weak scale, and further complicating the precise relation between the underlying parameters and $\langle H\rangle$ and $\langle \overline{H}\rangle$. The VEV for S is crucial, because combined with the dynamical super-potential, it provides supersymmetric masses for the fermionic components[4] of ψ, $\overline{\psi}$, χ and $\overline{\chi}$. Thus, we expect that in generic points in the parameter space, subject to quite mild constraints, phenomenologically viable EWSB and weak scale masses for the uncolored exotics result.

2.3 Light fermion masses

We have seen that the top Yukawa coupling and Higgs quartic are generated by the strong dynamics, and are naturally large. The remainder of the fermion masses can also be generated in the following ways.

2.3.1 Charged leptons

The lepton sector is entirely fundamental, so the required operators are dimension 5 at the preon level, to connect L_i, e_j and the composite Higgs $\overline{H}$. The needed underlying interactions are generated by integrating out the spectators H' and $\overline{H}'$ (just as in the original FH), and result in,

$$\begin{aligned} W_L &= y_{H'} H' P_1 \overline{P}_2 + y^e_{ij} \overline{H}' L_i e_j \\ &\to \left(\frac{y^e_{ij} y_{H'}}{4\pi} \frac{\Lambda}{M_H} \right) \overline{H} L_i e_j \,. \end{aligned} \tag{2.19}$$

As in the fat Higgs case, this is suppressed by Λ/M_H. However, a wide range of parameters is permitted given the smallness of the observed charged lepton masses.

2.3.2 Down-type quarks

The couplings between the fundamental left-handed quarks $Q_{1,2}$ to the fundamental right-handed down quarks, $d_{1,2,3}$ is also a dimension five operator. It can also be generated by the spectator Higgses,

$$\begin{aligned} W_{d1} &= y^d_{ij} \overline{H}' Q_i d_j \\ &\to \left(\frac{y_{H'} y^d_{ij}}{4\pi} \frac{\Lambda}{M_H} \right) \overline{H} Q_{1,2} d_i \,. \end{aligned} \tag{2.20}$$

[4]Alternately, one may introduce further spectators to marry ψ, $\overline{\psi}$, χ, and $\overline{\chi}$ with masses of order Λ through non-renormalizable operators mediated by a new set of spectator preons. While this results in a more minimal particle content below Λ (and reproduces precisely the FH scalar potential), it requires many more ingredients, and thus we prefer to accept the extra light states at the weak scale.

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We also need couplings between Q_3 and d_i, in order to have a bottom quark mass. This requires a dimension 6 interaction between preons, to connect Q_3 and $\overline{H}$ (both mesons) to d_i. This can be arranged by integrating out both P' and H', through the interactions,

$$\begin{aligned} W_{d2} &= y_{\overline{H}'} \overline{H}' \overline{P}_2 P' + y_{d_j} \overline{P}' P_3 d_j \\ &\to \left(y_{H'} y_{\overline{H}'} y_{d_j} \frac{\Lambda^2}{M_P M_H} \right) \overline{H} Q_3 d_i \,. \end{aligned} \tag{2.21}$$

Note that the NDA estimates do not include a 4π suppression in this case, which might point to bottom being naturally heavier than down or strange. At this point, the down-type quark mass matrix is generic - it contains no necessarily zero or very small entries. Thus, it is able to generate all of the down-type masses, and (after we generate the up and charm quark masses, below) is sufficient to generate the full CKM structure of the Standard Model.

2.3.3 Up-type quarks

Finally, we need a mass for the up and charm quarks, the top quark having already been arranged through the dynamical super-potential. Since the CKM mixing has already been arranged in the down-type sector, we do not pursue masses linking Q_3 with $u_{1,2}$ (or $Q_{1,2}$ with t_R) but instead just masses connecting Q_i with u_j where $i, j = 1, 2$. These can be generated by integrating out both P' and H',

$$\begin{aligned} W_u &= y^u_{ij} H' Q_i u_j + y_{P_1} \overline{P}' \overline{P}_1 \overline{P}_{\tilde{1}} \\ &\to \left(\frac{y^u_{ij} y_{\overline{H}'} y_{P_1}}{4\pi} \frac{\Lambda^2}{M_P M_H} \right) H Q_{1,2} u_{1,2} \,. \end{aligned} \tag{2.22}$$

And thus all Yukawas can be built by integrating out the spectator preons P' and Higgses H'.

2.3.4 Residual interactions

In addition to the light fermion Yukawa interactions described above, there are residual effects from integrating out the spectators H' and $\overline{H}'$. The first is that these massive fields mediate flavor-violating interactions of the form,

$$W_{\not{F}} = \left(\frac{y^u_{ij} y^d_{kl}}{M_H} \right) Q_i u_j Q_k d_l + \left(\frac{y^u_{ij} y^e_{kl}}{M_H} \right) Q_i u_j L_k e_l \,. \tag{2.23}$$

While not a consequence of the composite sector in our model, these types of interactions are often referred to as "compositeness operators" [16]. They lead to interactions involving two SM fermions and two of their scalar superpartners, and thus to anomalous flavor violation at the loop level. Given the large value of $M_H \gtrsim \Lambda \sim 1000\,\text{TeV}$, they are not expected to be in contradiction with data, though they are in the region where improved precision in future experiments could potentially see some of their effects.

The second operator is an induced μ-term for the composite EWSB Higgses H and $\overline{H}$,

$$W_\mu = \left(y_{P_1} y_{\overline{H}'} y_{H'} \frac{\Lambda^3}{M_H M_P} \right) \; H\overline{H} \equiv \mu \; H\overline{H} \,. \tag{2.24}$$

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As we saw above, a large μ term would lead to EW fine-tuning, and so we assume that the Yukawa interactions and/or the suppression from Λ/M_H is sufficient to bring this operator down to the weak scale.

Both of these features are a consequence of our having taken a minimal approach to the question of flavor, and not an "over-kill" approach as proposed in [17]. There is no problem to incorporate the over-kill framework in our SU(3) model, though since the contributions are not sizable enough to be dangerous, we choose to present the simpler and potentially more phenomenologically interesting case here.

2.4 Exotic quark masses

We have already seen that the VEV for the singlet S generates weak scale masses for the ψ and χ superfields for fairly generic parameters. We also need masses for the exotic quarks q_1, q_2, in order to avoid having these them appear at low energies. We introduce fundamental fields $\overline{q}_{(1,2)}$ to marry these exotics through the super-potential,

$$\begin{aligned} W_q &= y_{q_1}\overline{q}_1 P_3 \overline{P}_1 + y_{q_2}\overline{q}_2 P_3 \overline{P}_{\tilde{1}} \\ &\to \left(\frac{y_{q_1}}{4\pi}\Lambda\right)\overline{q}_1 q_1 + \left(\frac{y_{q_2}}{4\pi}\Lambda\right)\overline{q}_2 q_2\,, \end{aligned} \tag{2.25}$$

where we continue to include the NDA 4π estimates. Thus, we typically expect that q_1 and q_2 are the heaviest of the exotics.

3. Unification

One of the hallmark successes of the MSSM is the prediction of the unification of the gauge couplings. In this section we demonstrate that this success can also be preserved in our SU(3) FH model. Unlike the generations of the MSSM, our preons do not fill out complete SU(5) representations, and so it is clear that the standard structural successes of four dimensional GUTs is not present. However, it may be that unification of couplings results from "string unification" or from a higher dimensional theory with orbifold GUT breaking [20], in which case matter need not fill out complete representations.

The evolution of the gauge couplings takes place in two steps. Below the strong coupling scale Λ the matter content is that of the MSSM, including the composite Higgses and top quark, plus the weak scale exotics S, ψ, $\overline{\psi}$, χ, and $\overline{\chi}$. The fields S, χ, and $\overline{\chi}$ are singlets under the MSSM gauge groups, and thus do not contribute to the evolution of couplings at one loop. Thus, the couplings evolve as,

$$\frac{dg_i}{dt} = \beta_i \frac{g_i^3}{16\pi^2} \tag{3.1}$$

with

$$\beta_i = \left(-3, 1, \frac{39}{5}\right) \tag{3.2}$$

for $(\mathrm{SU}(3)_C, \mathrm{SU}(2)_W, \mathrm{U}(1)_Y)$, and we have normalized the hypercharge coupling in the usual SU(5) way, $\beta_1 = 3/5\beta_Y$.

Above the scale Λ the evolution includes the extra composites q_1 and q_2 (and their partners). More correctly, one should consider the evolution in terms of the preons as the relevant degrees of freedom at large scales, but the two descriptions are equivalent because of holomorphicity. In order to recover unification of couplings, we also include two vector-like pairs of spectator "unifons" which do not participate in the strong dynamics, and are doublets under $\mathrm{SU}(2)_W$ with no hyper-charge. Thus, above Λ we have,

$$\beta_i = (-2, 3, 9)\,, \tag{3.3}$$

and combining these together with $\Lambda \sim 1000\,\mathrm{TeV}$, we find unification of couplings at the level of 5% at a scale of $3 \times 10^{14}\,\mathrm{GeV}$. Such a low scale of unification could be problematic with respect to proton stability, but since there is no clear GUT structure the usual proton decay mediated by X, Y GUT bosons may not be present and could be further evaded by imposing some type of baryonic symmetry.

One might worry that the additional strong dynamics will spoil any true prediction of unification because of the extra strong dynamics threshold at Λ. In a supersymmetric theory, this is not a problem because the holomorphicity of the super-potential demands that the low energy couplings are determined only by the bare masses of the heavy fields [21]. Thus, our SU(3) FH theory has true unification at a level comparable to the MSSM.

4. Phenomenology

This model has some distinctive phenomenology, which helps to distinguish it from other supersymmetric theories. The MSSM super-partner phenomenology depends (as usual) quite crucially on the mechanism by which SUSY breaking is communicated to the MSSM fields, and thus is model-dependent. In order to avoid EW fine-tuning, it is important that the scalar partners of top be no more than a few hundred GeV. This requirement, combined with a model of SUSY breaking at high scales will also favor a gluino mass in this region (see [19] for models designed to evade this requirement). Stop masses of up to about 200 GeV (depending on decay mode and other super-partner masses) can be found in a variety of decay modes at the Tevatron [18], which can also typically discover gluinos provided their mass is less than 400 GeV [22]. The LHC is expected to be sensitive to gluino masses up to about 2 TeV [24].

4.1 Higgs

Including the S superfield, our theory has the additional singlet Higgs (containing additional neutral scalars and pseudo-scalars) which mixes through EWSB with the usual MSSM Higgses. This rich spectrum corresponds to various cases of the next-to-minimal supersymmetric standard model, and has been studied in great detail [23]. The mixing with the extra scalar state can lead to reduced Z-Z-h^0 and W-W-h^0 couplings, thus weakening the LEP II direct search limits. The fermionic component of S will also mix with the MSSM neutralinos, leading to a modification of the MSSM neutralino properties [25].

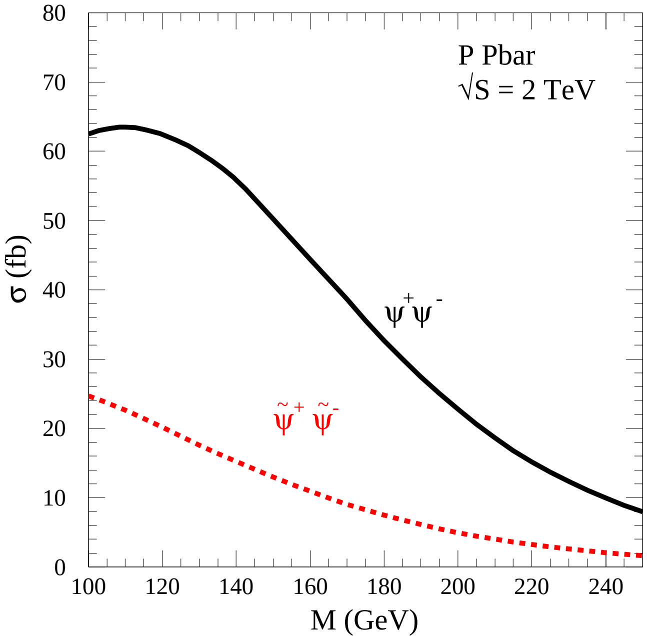

Figure 3: The cross sections for producing $\psi^+\psi^-$ and $\widetilde{\psi}^*\widetilde{\psi} + \widetilde{\overline{\psi}}^*\widetilde{\overline{\psi}}$ at the Tevatron.

The Higgs responsible for EWSB is generally quite a bit heavier than in the usual MSSM, because of the large value of λ_H which contributes to the Higgs mass. For large m_A, $\tan\beta \sim 1$ and $\Lambda \sim 1000\,\mathrm{TeV}$, the mass is expected to be around $140\,\mathrm{GeV}$, which is considerably higher than any reasonable value in the MSSM, and high enough that decays such as $H \to WW^*$ will begin to dominate. More exotic decay modes such as $H \to A^0A^0$ may occur, and can be very challenging for LHC Higgs searches [26]. In addition, large values of λ_H can lead to the charged Higgs being the lightest one, something that never occurs in the MSSM [6, 7].

4.2 Exotics

The model also has a number of additional chiral multiplets. The colored quark singlets q_1 and q_2 have masses of order Λ (and thus will probably not be produced at near future colliders), whereas the color neutral particles are expected to have masses $\lambda_\psi\langle S\rangle$, of order $v \sim 200\,\mathrm{GeV}$. We expect the lightest of these to be the singlet χ fields, and the charge ± 1 fields ψ should be slightly heavier, because of its non-zero hyper-charge. We expect that the scalar components will be slightly heavier than their fermionic partners because of SUSY-breaking contributions to the scalar masses.

The dynamically generated super-potential has a $\mathbb{Z}_2$ symmetry which has all of the exotic particles coupling in pairs. This symmetry could be imposed exactly, but more likely will be broken by interactions such as $\overline{q}_1 d_i d_j$, which allows the scalar $\overline{q}_1$ to decay directly into down-type quarks (or the fermionic q_1 to decay into two quarks and a gaugino). Since all of the exotic states must decay through $\overline{q}_1$ whose mass is of order $1000\,\mathrm{TeV}$, the exotics are typically very long lived and have complicated multi-particle final states. In the case

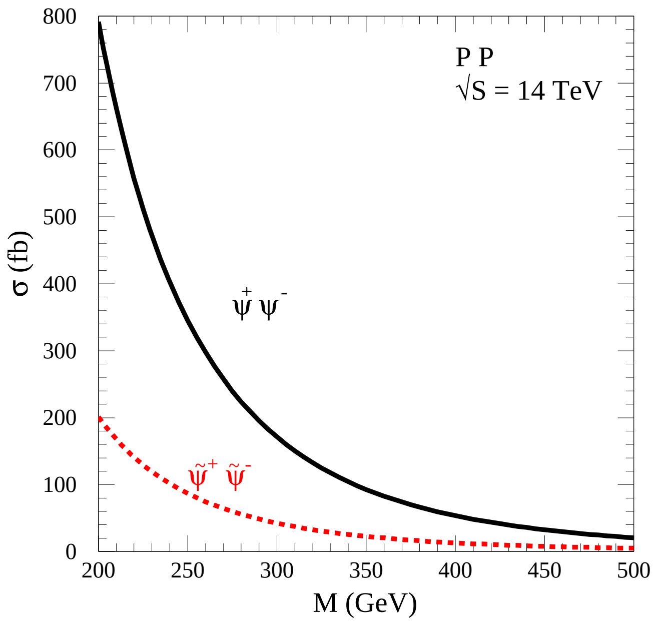

Figure 4: The cross sections for producing $\psi^+\psi^-$ and $\widetilde{\psi}^*\widetilde{\psi} + \widetilde{\overline{\psi}}^*\widetilde{\overline{\psi}}$ at the LHC.

of ψ, this results in electrically charged fermions and their scalar partners which are stable on length scales of the order of the detector, and thus appear as massive charged objects. Studies in ref. [27] considered such objects in the context of certain gauge-mediated SUSY breaking models and conclude that the Tevatron can discover them with $2\,\mathrm{fb}^{-}1$ at the 5σ level provided the production cross section is larger than about 100(10) fb for masses of 100(250) GeV. In figure 3 we plot the production cross sections for both the fermion (ψ) and scalars ($\widetilde{\psi}$ and $\widetilde{\overline{\psi}}$) at the Tevatron [28], through the partonic processes $q\overline{q} \to \gamma, Z \to \psi^+\psi^-$, and so forth for the scalars. Note that the scalar cross sections are suppressed relative to the fermionic ones because of the intermediate vector boson, which requires that the scalars be produced in the p-wave to conserve angular momentum. For a wide variety of masses, the Tevatron should be able to probe this scenario with $2\,\mathrm{fb}^{-1}$ of collected luminosity. The LHC should be able to produce and detect the charged quasi-stable particles up to even larger masses. The cross sections at the LHC are plotted in figure 4 [28], and it is expected that the LHC will cover the entire parameter space [29]. The χ and $\overline{\chi}$ particles will be produced much less copiously, and being electrically neutral and quasi-stable are very difficult to detect.

5. Conclusions

The fat Higgs is a fascinating alternative to the minimal supersymmetric standard model, which may naturally explain why LEP II did not discover the light CP even Higgs responsible for EWSB. In this article, we have examined an alternative to the minimal model based on an s-confining (at $\sim 1000\,\mathrm{TeV}$) SU(3) group which generates not only the MSSM Higgses and a singlet, but also the top quark as composites in the low energy theory. This naturally generates the large top Yukawa coupling as a residual of the strong dynamics, perhaps explaining why top is so much more massive than any other fermion of the Standard Model.

We are able to generate all of the observed flavor structure of the standard model, and predict that the Higgs mass and top mass are correlated because of the common origin of both couplings from the dynamical super-potential. This relieves some fine-tuning in the original FH model, and perhaps motivates the large top mass. Electroweak symmetry breaking happens in a way which is reminiscent of the FH, and does impose some mild conditions on the soft masses of the MSSM-like and exotic Higgses.

The model is compatible with unification of couplings, and results in some weak scale exotic states not seen the in the MSSM. These include quasi-stable electrically charged (± 1) objects for which there are good discovery prospects at the Tevatron run II once $2\,\mathrm{fb}^{-1}$ of data has been collected. These provide a means to distinguish this model from other supersymmetric theories, including the original fat Higgs itself. There are also interesting modifications to Higgs physics, with the most important one being the fact that the lightest CP even Higgs will typically be heavier than in the MSSM, even at tree level. Clearly, supersymmetric theories are likely to be richer than even the minimal models, and the next generation of colliders is likely to have an exciting time unrevealing the physics at the TeV scale.

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References

[1] ALEPH collaboration, D. Abbaneo et al., *A combination of preliminary electroweak measurements and constraints on the standard model*, hep-ex/0112021.

[2] Y. Okada, M. Yamaguchi and T. Yanagida, *Upper bound of the lightest Higgs boson mass in the minimal supersymmetric standard model*, *Prog. Theor. Phys.* **85** (1991) 1;
M. Carena, M. Quiros and C.E.M. Wagner, *Effective potential methods and the Higgs mass spectrum in the MSSM*, *Nucl. Phys.* **B 461** (1996) 407 [hep-ph/9508343];
H.E. Haber, R. Hempfling and A.H. Hoang, *Approximating the radiatively corrected Higgs mass in the minimal supersymmetric model*, *Z. Physik* **C 75** (1997) 539 [hep-ph/9609331];
J.R. Espinosa and I. Navarro, *Radiative corrections to the Higgs boson mass for a hierarchical stop spectrum*, *Nucl. Phys.* **B 615** (2001) 82 [hep-ph/0104047];
K. Tobe and J.D. Wells, *Higgs boson mass limits in perturbative unification theories*, *Phys. Rev.* **D 66** (2002) 013010 [hep-ph/0204196].

[3] M. Carena, J.R. Ellis, S. Mrenna, A. Pilaftsis and C.E.M. Wagner, *Collider probes of the MSSM Higgs sector with explicit CP-violation*, *Nucl. Phys.* **B 659** (2003) 145 [hep-ph/0211467].

[4] J.R. Ellis, J.F. Gunion, H.E. Haber, L. Roszkowski and F. Zwirner, *Higgs bosons in a nonminimal supersymmetric model*, *Phys. Rev.* **D 39** (1989) 844;
H.E. Haber and M. Sher, *Higgs mass bound in E_6 based supersymmetric theories*, *Phys. Rev.* **D 35** (1987) 2206;
M. Drees, *Comment on 'higgs boson mass bound in E_6 based supersymmetric theories.'*, *Phys. Rev.* **D 35** (1987) 2910;
K.S. Babu, X.G. He and E. Ma, *New supersymmetric left-right gauge model: higgs boson structure and neutral current analysis*, *Phys. Rev.* **D 36** (1987) 878;
M. Cvetič, D.A. Demir, J.R. Espinosa, L.L. Everett and P. Langacker, *Electroweak breaking and the mu problem in supergravity models with an additional* U(1), *Phys. Rev.* **D 56** (1997) 2861 [hep-ph/9703317].

[5] P. Batra, A. Delgado, D.E. Kaplan and T.M.P. Tait, *The Higgs mass bound in gauge extensions of the minimal supersymmetric standard model*, *JHEP* **02** (2004) 043 [hep-ph/0309149].

[6] P. Batra, A. Delgado, D.E. Kaplan and T.M.P. Tait, *Running into new territory in SUSY parameter space*, *JHEP* **06** (2004) 032 [hep-ph/0404251].

[7] R. Harnik, G.D. Kribs, D.T. Larson and H. Murayama, *The minimal supersymmetric fat Higgs model*, *Phys. Rev.* **D 70** (2004) 015002 [hep-ph/0311349].

[8] S. Chang, C. Kilic and R. Mahbubani, *The new fat Higgs: slimmer and more attractive*, *Phys. Rev.* **D 71** (2005) 015003 [hep-ph/0405267].

[9] J.A. Casas, J.R. Espinosa and I. Hidalgo, *The MSSM fine tuning problem: a way out*, *JHEP* **01** (2004) 008 [hep-ph/0310137];
A. Birkedal, Z. Chacko and Y. Nomura, *Relaxing the upper bound on the mass of the lightest supersymmetric Higgs boson*, *Phys. Rev.* **D 71** (2005) 015006 [hep-ph/0408329];
A. Birkedal, Z. Chacko and M.K. Gaillard, *Little supersymmetry and the supersymmetric little hierarchy problem*, *JHEP* **10** (2004) 036 [hep-ph/0404197];
R. Martinez, N. Poveda and J.-A. Rodriguez, *Upper bound of the lightest Higgs boson in a supersymmetric* $\mathrm{SU}(3)_L \times \mathrm{U}(1)$ *gauge model*, *Phys. Rev.* **D 69** (2004) 075013;
K.S. Babu, I. Gogoladze and C. Kolda, *Perturbative unification and Higgs boson mass bounds*, hep-ph/0410085.

[10] A. Maloney, A. Pierce and J.G. Wacker, *D-terms, unification and the Higgs mass*, hep-ph/0409127.

[11] J.R. Espinosa and M. Quiros, *On Higgs boson masses in nonminimal supersymmetric standard models*, *Phys. Lett.* **B 279** (1992) 92; *Upper bounds on the lightest Higgs boson mass in general supersymmetric standard models*, *Phys. Lett.* **B 302** (1993) 51 [hep-ph/9212305]; *Gauge unification and the supersymmetric light Higgs mass*, *Phys. Rev. Lett.* **81** (1998) 516 [hep-ph/9804235];
K. Tobe and J.D. Wells, *Higgs boson mass limits in perturbative unification theories*, *Phys. Rev.* **D 66** (2002) 013010 [hep-ph/0204196].

[12] A.G. Cohen, D.B. Kaplan and A.E. Nelson, *Counting 4pi's in strongly coupled supersymmetry*, *Phys. Lett.* **B 412** (1997) 301 [hep-ph/9706275].

[13] For pedagogical reviews, see, K.A. Intriligator and N. Seiberg, *Lectures on supersymmetric gauge theories and electric- magnetic duality*, *Nucl. Phys.* **45BC** *(Proc. Suppl.)* (1996) 1 [hep-th/9509066];
M.E. Peskin, *Duality in supersymmetric Yang-Mills theory*, hep-th/9702094;
J. Terning, *Non-perturbative supersymmetry*, hep-th/0306119.

[14] C. Csáki, M. Schmaltz and W. Skiba, *Confinement in $N = 1$ SUSY gauge theories and model building tools*, *Phys. Rev.* **D 55** (1997) 7840 [hep-th/9612207].

[15] CDF AND D0 collaborations, C.S. Hill *Top quark mass measurements*, FERMILAB-CONF-04-395-E, prepared for *5th Rencontres du Vietnam on particle physics and astrophysics: new views in particle physics* Vietnam 2004, Hanoi, Vietnam, 5–11 Aug 2004, SPIRES entry.

[16] E. Eichten, K.D. Lane and M.E. Peskin, *New tests for quark and lepton substructure*, *Phys. Rev. Lett.* **50** (1983) 811.

[17] H. Murayama, *Technicolorful supersymmetry*, hep-ph/0307293.

[18] R. Demina, J.D. Lykken, K.T. Matchev and A. Nomerotski, *Stop and sbottom searches in run II of the FermiLab Tevatron*, *Phys. Rev.* **D 62** (2000) 035011 [hep-ph/9910275].

[19] T. Kobayashi, H. Nakano and H. Terao, *Induced top Yukawa coupling and suppressed Higgs mass parameters*, *Phys. Rev.* **D 71** (2005) 115009 [hep-ph/0502006].

[20] Y. Kawamura, *Gauge symmetry reduction from the extra space $S_1/\mathbb{Z}_2$*, *Prog. Theor. Phys.* **103** (2000) 613 [hep-ph/9902423]; *Triplet-doublet splitting, proton stability and extra dimension*, *Prog. Theor. Phys.* **105** (2001) 999 [hep-ph/0012125];
L.J. Hall and Y. Nomura, *Gauge unification in higher dimensions*, *Phys. Rev.* **D 64** (2001) 055003 [hep-ph/0103125];
A. Hebecker and J. March-Russell, *The structure of GUT breaking by orbifolding*, *Nucl. Phys.* **B 625** (2002) 128 [hep-ph/0107039];
L.J. Hall, H. Murayama and Y. Nomura, *Wilson lines and symmetry breaking on orbifolds*, *Nucl. Phys.* **B 645** (2002) 85 [hep-th/0107245].

[21] N. Arkani-Hamed and H. Murayama, *Renormalization group invariance of exact results in supersymmetric gauge theories*, *Phys. Rev.* **D 57** (1998) 6638 [hep-th/9705189].

[22] SUGRA WORKING GROUP collaboration, S. Abel et al., *Report of the sugra working group for run II of the Tevatron*, hep-ph/0003154.

[23] D.J. Miller, R. Nevzorov and P.M. Zerwas, *The Higgs sector of the next-to-minimal supersymmetric standard model*, *Nucl. Phys.* **B 681** (2004) 3 [hep-ph/0304049];
U. Ellwanger, J.F. Gunion, C. Hugonie and S. Moretti, *NMSSM Higgs discovery at the LHC*, hep-ph/0401228.

[24] ATLAS AND CMS collaborations, S. Asai, *Supersymmetry at LHC*, *Eur. Phys. J. Direct.* **C 4S1** (2002) 17.

[25] F. Franke and S. Hesselbach, *Production of singlino dominated neutralinos in extended supersymmetric models*, *Phys. Lett.* **B 526** (2002) 370 [hep-ph/0111285];
U. Ellwanger and C. Hugonie, *Topologies of the (M+1)SSM with a singlino LSP at LEP2*, *Eur. Phys. J.* **C 13** (2000) 681 [hep-ph/9812427].

JHEP07(2005)023

[26] B.A. Dobrescu, G. Landsberg and K.T. Matchev, *Higgs boson decays to CP-odd scalars at the Tevatron and beyond*, *Phys. Rev.* **D 63** (2001) 075003 [hep-ph/0005308];
U. Ellwanger, J.F. Gunion and C. Hugonie, *Difficult scenarios for NMSSM Higgs discovery at the LHC*, hep-ph/0503203.

[27] SUSY Working Group collaboration, R. Culbertson et al., *Low-scale and gauge-mediated supersymmetry breaking at the FermiLab Tevatron run II*, hep-ph/0008070.

[28] Adapted from the contribution by J.L. Hewett and T.M.P. Tait to: LHC/LC Study Group collaboration, G. Weiglein et al., *Physics interplay of the LHC and the ILC*, hep-ph/0410364.

[29] S. Ambrosanio et al., *SUSY long-lived massive particles: detection and physics at the LHC*, hep-ph/0012192.

JHEP07(2005)023