Superconformal Technicolor

Aleksandr Azatov* and Jamison Galloway†
Dipartimento di Fisica, Università di Roma "La Sapienza" and INFN Sezione di Roma, I-00185 Rome, Italy

Markus A. Luty‡
Physics Department, University of California Davis, Davis, California 95616, USA
(Received 27 September 2011; published 24 January 2012)

In supersymmetric theories with a strong conformal sector, soft supersymmetry breaking at the TeV scale naturally gives rise to confinement and chiral symmetry breaking at the same scale. We consider two such scenarios, one where the strong dynamics induces vacuum expectation values for elementary Higgs fields, and another where the strong dynamics is solely responsible for electroweak symmetry breaking. In both cases, the mass of the Higgs boson can exceed the LEP bound without tuning, solving the supersymmetry naturalness problem. A good precision electroweak fit can be obtained, and quark and lepton masses are generated without flavor-changing neutral currents. In addition to standard supersymmetry signals, these models predict production of multiple heavy standard model particles (t, W, Z, and b) from decays of resonances in the strong sector.

DOI: 10.1103/PhysRevLett.108.041802 PACS numbers: 12.60.Jv, 12.60.Nz, 14.80.Da

Introduction.—Supersymmetry (SUSY) is widely considered to be the most plausible framework for physics beyond the standard model of particle physics. It offers an elegant explanation of the fact that the electroweak breaking scale ~100 GeV is much smaller than the Planck scale $\sim 10^{19}$ GeV, without fine-tuning fundamental parameters. The minimal supersymmetric standard model (MSSM) also contains a viable dark matter candidate and gives a calculable framework for addressing other fundamental issues in particle physics and cosmology. However, there is a serious problem with electroweak symmetry breaking in the MSSM: the lightest Higgs boson has a mass that is generically $m_h < m_Z \simeq 90$ GeV, while the experimental bound from LEP is $m_h > 115$ GeV [1]. The Higgs boson mass can be raised at the cost of reintroducing tuning at the 1% level, or by extending the model in various ways [2]. In this Letter, we propose to solve this problem by combining supersymmetry with strong dynamics at the TeV scale. A companion paper [3] gives many additional details.

The electroweak scale in the MSSM is determined by the scale of soft SUSY breaking. We assume that in addition there is a strongly coupled sector of the theory with conformal (scale) invariance. An example of such a sector is SUSY QCD with $N_f \simeq 2N_c$ [4]. Soft SUSY breaking in the strong sector also softly breaks the conformal invariance. SUSY breaking in the strong sector gives mass to all scalars (since only unbroken SUSY can forbid these masses), while fermions generally remain massless due to unbroken chiral symmetries. It is therefore very plausible that the dynamics of SUSY QCD at the SUSY breaking scale is qualitatively similar to non-SUSY QCD, i.e., the theory confines and breaks chiral symmetry at this scale. Since the coupling is already strong at the SUSY breaking scale, these effects occur at this scale. In such models the strong sector can dynamically break electroweak symmetry, as in technicolor models [5]. Since the scale of dynamical electroweak symmetry breaking is determined by the soft breaking of conformal symmetry, this is a SUSY version of conformal technicolor [6], so we refer to it as superconformal technicolor [7]. We assume that the SUSY breaking scale is the same order of magnitude in the MSSM and the strong sector, which is natural in many models of SUSY breaking. This class of models therefore gives a plausible framework for SUSY *and* strong dynamics at the same scale. In Ref. [8] this mechanism was employed with a SUSY breaking scale above the electroweak scale to give a realistic model for flavor in conformal technicolor (the pioneering work in this direction is Ref. [9]). In the present work, we investigate SUSY breaking and strong dynamics at the TeV scale. Early attempts in this direction posited dynamical SUSY breaking at the TeV scale [10], but this is problematic for both theoretical and phenomenological reasons. A realistic model combining SUSY and strong dynamics at the TeV scale was constructed in Ref. [11]. The present work improves on that work in giving a general and robust mechanism for the coincidence of the scales of SUSY breaking and strong dynamics.

Induced electroweak symmetry breaking.—In these models there are two potential sources of electroweak symmetry breaking, the strong sector and the elementary Higgs fields of the MSSM. We first consider a scenario where electroweak symmetry breaking is induced by the strong sector, but the W and Z masses are dominated by the contribution from the elementary Higgs fields. A minimal strong sector has fields transforming under $SU(2)_{\rm SC} \times SU(2)_W \times U(1)_Y$ as

$$\Psi \sim (2,2)_0, \qquad \tilde{\Psi}_1 \sim (2,1)_{(1/2)}, \qquad \tilde{\Psi}_2 \sim (2,1)_{-(1/2)}, \quad (1)$$

plus 2 copies of $(2,1)_{(1/2)} \oplus (2,1)_{-(1/2)}$ fields that play no role in breaking electroweak symmetry. (The hypercharge assignments ensure that there are no fractionally charged states in the strong sector.) The fields Ψ and $\tilde{\Psi}$ have the quantum numbers of the technifermions of minimal technicolor [5]. The soft SUSY breaking terms explicitly break the global symmetry of the strong sector to $SU(2)_L \times SU(2)_R$. SUSY breaking in the strong sector is assumed to trigger confinement and chiral symmetry breaking by a fermion condensate $\langle \Psi\tilde{\Psi}\rangle \neq 0$, as in technicolor. (It is also natural to have a larger group of approximate symmetries due to the special structure of the soft SUSY breaking terms, in which case there will be additional light pseudo-Nambu Goldstone bosons.) Stabilizing runaway directions in the strong sector requires additional interactions, which are discussed in Ref. [3]. The strong sector is coupled to the MSSM Higgs fields via the superpotential couplings

$$W = \lambda_u H_u (\Psi\tilde{\Psi}_2) + \lambda_d H_d (\Psi\tilde{\Psi}_1). \quad (2)$$

The operators $\Psi\tilde{\Psi}$ have dimension $\simeq \frac{3}{2}$ above the SUSY breaking scale, and so the couplings $\lambda_{u,d}$ have mass dimension $\simeq +\frac{1}{2}$. We require that the couplings $\lambda_{u,d}$ be large enough to be important at the SUSY breaking scale, but not nonperturbatively large. This is a coincidence of scales between a relevant SUSY preserving coupling and the SUSY breaking scale, similar to the problem of why the superpotential term $\mu H_u H_d$ has a coupling $\mu \sim 100$ GeV. The simplest solution to the "μ problem" is the Giudice-Masiero mechanism [12], and Ref. [3] gives a generalization of this mechanism that can explain the required values of $\lambda_{u,d}$.

This solution of the coincidence problem can solve another potential problem for this class of models. The renormalization of SUSY breaking terms in the strong sector suppresses all scalar masses for the strong fields except those proportional to global symmetry generators [13]. This necessarily results in negative mass-squared terms for some of the strongly coupled fields and an unstable vacuum. This problem can be solved by adding additional elementary fields coupled to the strong fields in the same manner as the Higgs coupling, Eq. (2). These can lift the flat directions provided that these couplings are the same order of magnitude as the SUSY breaking terms. A detailed model is described in Ref. [3].

We assume that the strong sector dynamically breaks electroweak symmetry with order parameter f somewhat below what is required to explain the W and Z masses, e.g., $f \simeq 100$ GeV. We expect that the strong sector contains massive "hadron" states at a scale $\Lambda \sim 4\pi f \sim$ TeV. We assume that the elementary Higgs fields $H_{u,d}$ have masses below Λ, so the effective theory below the scale Λ contains these fields. The $SU(2)_L \times SU(2)_R$ symmetry is nonlinearly realized in the low-energy effective theory by $\Sigma(x) \in SU(2)$ transforming as $\Sigma \mapsto L\Sigma R^\dagger$. The elementary Higgs fields and couplings in Eq. (2) can be combined into a spurion

$$\mathcal{H}\lambda = \begin{pmatrix} \lambda_d H_d^0 & \lambda_u H_u^- \\ \lambda_d H_d^+ & \lambda_u H_u^0 \end{pmatrix} \mapsto L\mathcal{H}\lambda R^\dagger. \quad (3)$$

The strong dynamics generates new contributions to the Higgs potential

$$\Delta V_{\text{eff}} = \frac{\Lambda^4}{16\pi^2}\left[\frac{c_1}{\Lambda}\,\text{tr}(\Sigma^\dagger \mathcal{H}\lambda) + \text{H.c.} + \mathcal{O}\left(\frac{\mathcal{H}\lambda}{\Lambda}\right)^2\right] \quad (4)$$

with $c_1 \sim 1$ [14]. This contains a linear term for the Higgs fields, so the Higgs fields get vacuum expectation values (VEVs) even for $m_{H_{u,d}}^2 > 0$, which we assume to be the case. (In standard SUSY scenarios $m_{H_{u,d}}^2 > 0$ at high scales and renormalization group running results in $m_{H_u}^2 < 0$ at the TeV scale, but more general boundary conditions at high scales can lead to $m_{H_{u,d}}^2 > 0$ at the TeV scale.) We assume that this generates VEVs for the elementary Higgs fields with $v_u^2 + v_d^2 \gg f^2$ (e.g., for $f \simeq 100$ GeV, $\sqrt{v_u^2 + v_d^2} \simeq 225$ GeV). The higher order terms in Eq. (4) are negligible if $m_{H_{u,d}} \ll 4\pi f^2/v$. Note that we get a stable minimum even if we neglect the quartic terms in the potential and $B\mu$ term, and the physical Higgs masses are given by the quadratic terms in the potential in this limit. Including the full potential, the predictions are more complicated, but the Higgs masses are still arbitrary parameters depending on the SUSY breaking mass terms. The physical Higgs masses can have any value ~100 GeV without fine-tuning, so this completely solves the SUSY Higgs mass problem.

The quark and lepton masses arise from conventional Yukawa couplings to $H_{u,d}$, which have a minimal flavor-violating structure. Since $\langle H_{u,d}\rangle$ is the dominant source of electroweak symmetry breaking, the Yukawa couplings are perturbative, even for the top quark. Therefore, there is no flavor problem associated with the strong dynamics.

We now turn to the phenomenology of this model. Early work on technicolor theories with Higgs scalars can be found in Refs. [15]. We first discuss the precision electroweak fit. The strong sector has $N_c = 2$ and only one weak doublet, so the contributions to the S and T parameters from the strong sector are not dangerously large to begin with, and there are large theoretical uncertainties in their values. In fact, general theoretical arguments suggest that the S parameter is suppressed in theories that are conformal above the chiral symmetry breaking scale [16]. Recent lattice simulations give some support for this behavior [17]. In the present model the IR contribution to S from the strong sector is reduced compared to a conventional technicolor theory because the pseudo-Nambu-Goldstone bosons (PNGBs) are heavy, and because there is a light Higgs field in the spectrum. Custodial symmetry can be broken in the strong sector by $\lambda_u v_u \neq \lambda_d v_d$. We assume that this contribution to the T parameter is positive, as suggested by perturbation theory. This means that the theory has an adjustable parameter that allows a good precision electroweak fit (similar to the Higgs mass in the standard model). We can easily obtain a good precision electroweak fit, even if we assume (pessimistically) that the UV contribution to the S parameter has the value obtained by extrapolation from QCD [18]. This is illustrated in Fig. 1.

Another important precision electroweak constraint is the coupling of the Z to left-handed b quarks. In this model the leading correction enters at $\mathcal{O}(y^2\lambda^2)$, with y the Yukawa coupling to standard model quark fields. The coupling $g_{Z\bar{b}b}$ agrees with the standard model at the 0.25% level, which gives a constraint $v < 5.6f$. This is easily satisfied given the other constraints we have considered, and we conclude that this coupling does not significantly restrict the viable parameter space.

We now discuss the signals at the LHC. In addition to the standard MSSM signatures, the theory has new signatures from the extended Higgs sector. In the simplified limit discussed above, the CP even scalars have masses $m_{H_{u,d}}$ while for the CP odd scalars we have (for $f \ll v$)

$$m^2_{A_1^0} = m_{H_1^\pm} = \frac{m^2_{H_u} m^2_{H_d}}{m^2_{H_u}\sin^2\beta + m^2_{H_d}\cos^2\beta}, \tag{5}$$

$$m^2_{A_2^0} = m_{H_2^\pm} = \frac{v^2}{f^2}(m^2_{H_u}\sin^2\beta + m^2_{H_d}\cos^2\beta), \tag{6}$$

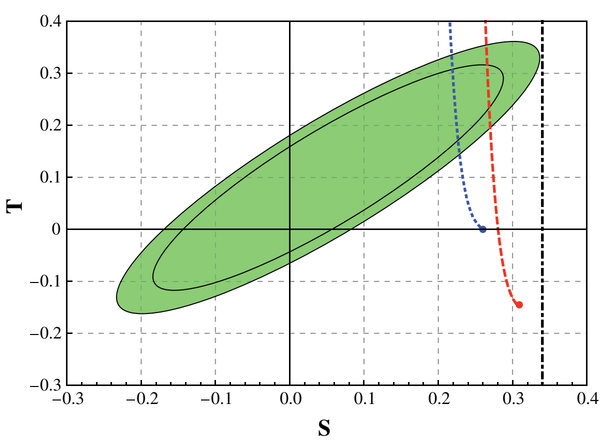

FIG. 1 (color online). Precision electroweak fit. The inner (outer) ellipse is the 95% (99%) confidence level allowed region in the S, T plane with reference Higgs mass 120 GeV [19]. The dotted blue (dashed red) line corresponds to models with a light Higgs boson at 130 (350) GeV, with $f = 100$ GeV, $\tan\beta = 2$, and $B\mu = 0$. The lines end when $\lambda_u v_u \simeq \lambda_d v_d$, where T is dominated by the light Higgs contribution. The dot-dashed black line is for the model with no light Higgs boson. The plot assumes that the UV contribution to the S parameter is given by the QCD value, while the UV contribution to the T parameter is estimated using naive dimensional analysis.

where $\tan\beta = v_u/v_d$. The heavier mass eigenstates A_2^0 and $H_2^\pm$ are dominantly PNGBs from the strong sector, with mixing of order f/v with the elementary Higgs fields. The A_2^0 can be singly produced by gluon fusion via a top loop, with a rate suppressed by f^2/v^2. A_2^0 and $H_2^\mp$ can also be pair produced via heavy resonances in the strong sector. Dominant decay modes are $A_2^0 \to \bar{t}t$, $W^\pm H^\mp$, Zh^0, $A_1^0 h^0$ and $H_2^+ \to \bar{b}t$, W^+h^0, $H_1^+h^0$. The h^0 decays dominantly to $\bar{b}b$ or WW/ZZ depending on its mass, so this leads to events with multiple heavy standard model particles (W, Z, t and/or b). Another signal is resonances in the strong sector with masses of order $4\pi f \sim$ TeV. Analogy with QCD suggests that the theory may have a prominent isotriplet vector resonance, the ρ_T. This can be singly produced via mixing with the W and Z of order $g/4\pi$, or via weak boson fusion. The ρ_T will generally have strong decays to pairs of PNGBs, but because of the large elementary Higgs VEVs, the A_2^0 and $H_2^\pm$ masses can be sufficiently large that decays to these states are kinematically forbidden. The effective field theory expansion breaks down in this regime, but we still expect it to be qualitatively reliable. In this case the ρ_T will be a narrow resonance, similar to a W' and Z'. Techniscalars charged under $SU(2)_L$ and $SU(2)_R$ generally have different SUSY breaking masses, so there need not be any approximate symmetry that interchanges $SU(2)_L$ and $SU(2)_R$, analogous to parity in QCD. This means that ρ_T can decay to either WW or WWW. The ρ_T can also decay via mixing with the W and Z.

Strong electroweak symmetry breaking.—We now consider another scenario where there are no elementary Higgs fields below the TeV scale, and electroweak symmetry is broken entirely by the strong sector. This arises in a different parameter regime of the model described above, as follows. We assume that the couplings $\lambda_{u,d}$ in Eq. (2) get strong at a scale $\Lambda_* >$ TeV. Results on nonperturbative dynamics of SUSY gauge theories [4] indicate that below the scale Λ_* the theory flows to a new fixed point where these couplings are strong. In this new fixed point, $H_{u,d}$ become operators of the strong sector with dimension $\simeq \frac{3}{2}$. This means that the Yukawa couplings of $H_{u,d}$ to quarks and leptons become irrelevant interactions below the scale Λ_*, scaling as $(E/\Lambda_*)^{1/2}$. In order to avoid too much suppression for the top quark mass, we cannot have Λ_* arbitrarily far above the TeV scale. If $\Lambda_* \gg$ TeV the top quark Yukawa coupling gets strong at some scale above Λ_*, indicating top quark compositeness at high scales. Alternatively, models with $\Lambda_* \sim$ TeV are natural with a mechanism to explain the coincidence of scales, as described above. For $\Lambda_* \gtrsim$ TeV, quark and lepton masses arise from irrelevant interactions at the TeV scale, as in technicolor. However, these interactions originate from Yukawa couplings with minimal flavor violation, and there is no flavor problem associated with the strong breaking of electroweak symmetry.

At the TeV scale, soft SUSY breaking in the strong sector is assumed to trigger confinement and electroweak symmetry breaking, as discussed above. The soft SUSY breaking terms can be chosen so that the strong sector has a minimal symmetry breaking structure $SU(2)_L \times SU(2)_R \rightarrow SU(2)$, so the only strong degrees of freedom below the TeV scale are the longitudinal components of the W and Z. The spectrum at the TeV scale therefore includes all of the MSSM fields minus the Higgs sector, with strong resonances at the scale $4\pi v \sim 3$ TeV.

A good precision electroweak fit can be obtained with the help of a T parameter induced by $\lambda_u \neq \lambda_d$. Assuming that the S parameter is given by the QCD value, the precision electroweak fit is shown in Fig. 1. A good fit can be obtained if the UV contribution to the S parameter is reduced compared to this estimate (as we expect, as discussed above), and the contribution to the T parameter from $\lambda_u \neq \lambda_d$ is positive (as expected from perturbation theory). The correction to $y_{Z\bar{b}b}$ is of order 0.8%, with large theoretical uncertainties. This is roughly 3 times the experimental precision so there is some tension, but given the large uncertainties this does not rule out the model. The collider phenomenology consists of SUSY signals, plus technicolor resonances at the 3 TeV scale. The ρ_T can decay to both WW and WWW as described above, which distinguishes it from the conventional technirho.

Conclusions.—We have described models that solve the SUSY Higgs mass problem via strong dynamics at the TeV scale. The models consist of the MSSM plus a sector with a strong conformal fixed point. In such models, it is natural for the strong sector to dynamically break electroweak symmetry at the soft SUSY breaking scale. We considered two scenarios, one in which the strong breaking of electroweak symmetry induces the elementary Higgs VEVs, and one in which strong electroweak symmetry breaking dominates. In both scenarios the experimental bounds on light Higgs bosons are easily satisfied without tuning, and no additional flavor problem is introduced. Both scenarios have a dark matter candidate. However, gauge coupling unification is no longer a prediction of the minimal model described here, since the strong sector affects the evolution of the $SU(2)_W \times U(1)_Y$ gauge couplings but not $SU(3)_C$. Unification can be accommodated with additional matter fields, which however have no other apparent motivation in this framework. In conclusion, we believe that this is a plausible framework for electroweak symmetry breaking, and that the new signals suggested by these models deserve additional investigation.

We thank R. Contino, R. Kitano, T. Okui, and J. Terning for discussions. We also thank J. Serra for collaboration at early stages of this work. M. A. L. was supported by DOE Grant No. DE-FG02-91-ER40674.

*aleksandr.azatov@roma1.infn.it
†jamison.galloway@roma1.infn.it
‡luty@physics.ucdavis.edu

[1] R. Barate *et al.* (LEP Higgs Working Group), Phys. Lett. B **565**, 61 (2003).
[2] For an entry into the literature, see A. Birkedal, Z. Chacko, and Y. Nomura, Phys. Rev. D **71**, 015006 (2005); A. Maloney, A. Pierce, and J. G. Wacker, J. High Energy Phys. 06 (2006) 034; S. Chang, C. Kilic, and R. Mahbubani, Phys. Rev. D **71**, 015003 (2005); S. Chang, R. Dermisek, J. F. Gunion, and N. Weiner, Annu. Rev. Nucl. Part. Sci. **58**, 75 (2008).
[3] A. Azatov, J. Galloway, and M. A. Luty, Phys. Rev. D (to be published).
[4] N. Seiberg, Nucl. Phys. B **435**, 129 (1995).
[5] For a review, see C. T. Hill and E. H. Simmons, Phys. Rep. **381**, 235 (2003).
[6] M. A. Luty and T. Okui, J. High Energy Phys. 09 (2006) 070; M. A. Luty, J. High Energy Phys. 04 (2009) 050.
[7] This name has been previously used in for models that do not use the conformal technicolor mechanism to break electroweak symmetry in M. Antola, S. Di Chiara, F. Sannino, and K. Tuominen, arXiv:1001.2040; arXiv:1009.1624.
[8] J. A. Evans, J. Galloway, M. A. Luty, and R. A. Tacchi, J. High Energy Phys. 04 (2011) 003.
[9] S. Samuel, Nucl. Phys. B **347**, 625 (1990); M. Dine, A. Kagan, and S. Samuel, Phys. Lett. B **243**, 250 (1990).
[10] S. Dimopoulos and S. Raby, Nucl. Phys. B **192**, 353 (1981); M. Dine, W. Fischler, and M. Srednicki, Nucl. Phys. B **189**, 575 (1981).
[11] M. A. Luty, J. Terning, and A. K. Grant, Phys. Rev. D **63**, 075001 (2001).
[12] G. F. Giudice and A. Masiero, Phys. Lett. B **206**, 480 (1988).
[13] M. A. Luty and R. Rattazzi, J. High Energy Phys. 11 (1999) 001.
[14] H. Georgi, Phys. Lett. B **298**, 187 (1993).
[15] E. H. Simmons, Nucl. Phys. B **312**, 253 (1989); C. D. Carone and H. Georgi, Phys. Rev. D **49**, 1427 (1994).
[16] R. Sundrum and S. D. H. Hsu, Nucl. Phys. B **391**, 127 (1993).
[17] T. Appelquist *et al.* (LSD Collaboration), Phys. Rev. Lett. **106**, 231601 (2011).
[18] M. E. Peskin and T. Takeuchi, Phys. Rev. Lett. **65**, 964 (1990); B. Holdom and J. Terning, Phys. Lett. B **247**, 88 (1990).
[19] H. Flacher, M. Goebel, J. Haller, A. Hocker, K. Monig, and J. Stelzer, Eur. Phys. J. C **60**, 543 (2009).