Single Sector Supersymmetry Breaking

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We review recent work on realistic models that break supersymmetry dynamically and give rise to composite quarks and leptons, all in a single sector. These models have a completely natural suppression of flavor-changing neutral currents, and the hierarchy of Yukawa couplings is explained by the dimensionality of composite states. The generic signatures are unification of scalar masses with different quantum numbers at the compositeness scale, and lighter gaugino, Higgsino, and third-generation sfermion masses.

I. INTRODUCTION

One of the most exciting results of the recent progress in understanding strongly coupled supersymmetric gauge theories [1] is that it allows the exploration of new possibilities for the realization of supersymmetry (SUSY) in nature. The most important example is the use of dynamical SUSY breaking to explain the origin of the SUSY breaking scale. In recent years, large classes of SUSY gauge theories have been discovered that exhibit dynamical SUSY breaking through a variety of different mechanisms [2,3], and realistic models have been built using these theories as building blocks in both (super)gravity-mediated and gauge-mediated frameworks.

In these conventional approaches to SUSY model building, SUSY breaking arises in a separate sector consisting of fields that are neutral under the standard model (SM) gauge group, and SUSY breaking is communicated to the observable fields by messenger (gauge or gravitational) interactions. It is clearly important to know whether such a 'modular' structure is required in order for SUSY to be the solution of the hierarchy problem, or if simpler models without disjoint sectors are possible. In [4,5], realistic models were constructed that do not require a separate SUSY breaking sector. In these models, SUSY is broken dynamically by strongly coupled fields that are charged under the SM gauge group, giving rise to composites with the quantum numbers of quarks and leptons. The composite sfermions have SUSY breaking masses induced directly by the strong dynamics, while the masses of the composite fermions are protected by unbroken chiral symmetries. The masses of the elementary gauginos and sfermions arise from gauge mediation (via the composite scalars); they are therefore smaller than the composite scalar masses, which are necessarily in the range of 1–10 TeV.

If we make the simplest assumption that the first two generations are composite while the third is elementary, we automatically gain a partial understanding of the observed hierarchy of fermion masses. The reason is that all Yukawa couplings involving composite states must arise from higher-dimension operators in the fundamental theory, and are thus suppressed, while the top Yukawa coupling can be order one. A highly non-trivial feature of this scenario is that it does not lead to excessive flavor-changing neutral currents (FCNC's) from squark non-degeneracy even if the flavor sector has no flavor symmetry. This is because the strong composite dynamics is flavor-blind, and so the composite scalar masses are degenerate to high accuracy, with small corrections due to perturbative flavor-breaking couplings.

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Such scenarios for single sector SUSY breaking have several interesting generic phenomenological implications. First, as mentioned above, gaugino and stop masses will be much smaller than the masses of composite squarks and sleptons. Second, the composite scalar masses unify at the compositeness scale. Third, if we assume that the Yukawa interactions are generated by new physics at a flavor scale above the compositeness scale without special flavor symmetries, predictions for flavor-changing processes such as $\mu \to e\gamma$ are plausibly within experimental reach.

We have constructed large classes of supersymmetric gauge theories [5] with the non-perturbative dynamics required for this kind of model-building. At low energies the models either confine (like the models of [4]), have conformal fixed points, or are magnetically free. All of the models discussed here have only local SUSY breaking minima. Also, they are not calculable, and require dynamical assumptions to be phenomenologically viable. However, in many of the models the existence of a local SUSY breaking minimum with composite fermions can be established without dynamical assumptions. This shows that the combination of compositeness and SUSY breaking is not exotic, and suggests that further exploration of the connection between these phenomena is worthwhile.

II. MASS SCALES AND PHENOMENOLOGY

In this Section, we describe the most important qualitative features of single sector models. We will then focus on three example models: a 'meson' model where the first two generations correspond to dimension 2 operators; a 'dimensional hierarchy' model in which the first generation corresponds to dimension 3, the second to dimension 2, and the third generation is elementary (dimension 1); and a speculative model where the effective composite operator dimension is $\frac{3}{2}$. We want to emphasize the fact that the phenomenology is very rich, and is largely independent of the details of specific models. More thorough discussions can be found in Refs. [4,5].

A. SUSY Breaking and Compositeness

We first explain the mechanism that gives rise to SUSY breaking and compositeness. The known models have a strong gauge group of the form $G_{\rm comp} \times G_{\rm lift}$, where both groups are asymptotically free and $\Lambda_{\rm comp} \gg \Lambda_{\rm lift}$. The scale $\Lambda_{\rm comp}$ is the compositeness scale: composite quarks and leptons become strongly interacting at the scale $\Lambda_{\rm comp}$. Direct bounds on the compositeness scale imply that $\Lambda_{\rm comp} \gtrsim 2$ TeV. The role of the gauge group $G_{\rm lift}$ is to generate a dynamical superpotential that lifts the vacuum degeneracy and gives rise to a local SUSY breaking minimum.

The models contain the following fields

	$G_{\rm lift}$	$G_{\rm comp}$	$G_{\rm global}$
Q	$\square$	$\square$	$\mathbf{1}$
L	$\overline{\square}$	$\mathbf{1}$	$\overline{\square}$
$\bar{U}$	$\mathbf{1}$	$\overline{\square}$	$\square$
P	$\mathbf{1}$	R	$\mathbf{1}$

where the representation R may be highly reducible (implying additional global symmetries). In addition, the model has a tree-level superpotential

$$W = \lambda Q L \bar{U}. \tag{1}$$

$G_{\rm global}$ must be large enough that it contains $SU(3)_c \times SU(2)_L \times U(1)_Y$ as a subgroup. There are additional requirements on the model in order for this model to have a local SUSY breaking minimum. We choose $G_{\rm global}$ such that classically there is a flat direction with $\bar{U} \neq 0$ where Q and L are massive and $G_{\rm comp}$ is completely broken. Nonperturbative $G_{\rm lift}$ dynamics (gaugino condensation) lift this flat direction via a dynamical superpotential of the form

$$W_{\rm dyn} \sim \Lambda_{\rm lift}^{3-r} \bar{U}^r. \tag{2}$$

Whether this superpotential forces $\bar{U}$ to large or small values depends on the value of r, but it also depends on the effective Kähler potential for $\bar{U}$. For $\bar{U} \gg \Lambda_{\rm comp}$, $G_{\rm comp}$ is broken at weak coupling and the Kähler potential is smooth in $\bar{U}$, so the potential slopes toward the origin for $r > 1$. For $\bar{U} \ll \Lambda_{\rm comp}$ the $G_{\rm comp}$ dynamics changes the Kähler potential for $\bar{U}$. For example, if the $G_{\rm comp}$ dynamics is confining, the Kähler potential will be smooth in terms of a 'composite' field $B = (\bar{U}^n)$. The superpotential can then be written

$$W_{\rm dyn} \sim B^{r/n}, \tag{3}$$

which corresponds to a potential that slopes *away* from the origin if $r/n < 1$. Therefore, for $1 < r < n$ there is no SUSY minimum for any value of $\bar{U}$, and there is a SUSY breaking minimum near the border between the region of validity of the confined and Higgs descriptions. In general, there are other moduli corresponding to excitations of P that must be stabilized; in some models this requires an additional renormalizable term in the superpotential, while in other models it requires a dynamical assumption. The minimum occurs for

$$\langle \bar{U} \rangle \sim \frac{\sqrt{N} \Lambda_{\rm comp}}{4\pi}, \tag{4}$$

where N is the number of 'colors' of $G_{\rm comp}$.

This mechanism also occurs in cases where the $G_{\rm comp}$ dynamics gives rise to a conformal fixed point (in the limit where we turn off $G_{\rm lift}$), provided that the $\bar{U}$ anomalous dimensions are sufficiently large. As long as $\bar{U} \ll \Lambda_{\rm comp}$ the $G_{\rm comp}$ dynamics is controlled by the infrared fixed point. Recall that we are assuming that $G_{\rm lift}$ is weak at the scale $\Lambda_{\rm comp}$, so the non-perturbative superpotential can be viewed as a perturbation. The 1PI potential for $\bar{U}$ is therefore

$$V_{\rm 1PI} \simeq (K_{\rm 1PI}^{-1})_{\bar{U}^\dagger \bar{U}} \left| \frac{\partial W_{\rm dyn}}{\partial \bar{U}} \right|^2, \tag{5}$$

where $K_{\rm 1PI}$ is the 1PI Kähler potential evaluated at the conformal fixed point. The scaling dimension of the Kähler metric $(K_{\rm 1PI})_{\bar{U}^\dagger \bar{U}}$ is $2 - 2d_{\bar{U}}$, where $d_{\bar{U}}$ is the scaling dimension of $\bar{U}$. Therefore,

$$(K_{\rm 1PI}^{-1})_{\bar{U}^\dagger \bar{U}} \sim \bar{U}^{2(d_{\bar{U}}-1)/d_{\bar{U}}}. \tag{6}$$

This forces the potential to slope away from the origin for

$$\frac{1 - d_{\bar{U}}}{d_{\bar{U}}} > r - 1. \tag{7}$$

One might worry that this argument relies on a 'Higgs' description in terms of the elementary field $\bar{U}$ in a regime where the theory is strongly coupled. In many cases there is an alternate description in terms of a weakly-coupled dual theory. For example, if $G_{\rm comp} = SU(N)$ with F 'flavors', the theory has an infrared fixed point for $\frac{3}{2}N < F < 3N$ [6]. There is a dual description in terms of a theory with gauge group $SU(F - N)$ in which the 'baryon' operator $\bar{U}^N$ in the original theory is mapped to an operator $\bar{u}^{F-N}$ in the dual. For F near $\frac{3}{2}N$ the dual description is weakly coupled, and the considerations of the previous paragraph can be made rigorous. One finds that the behavior of the Kähler potential agrees precisely with Eq. (5). This equivalence between the 'Higgs' and 'dual' descriptions can be viewed as a generalization of the usual 'complementarity' [7] for theories with scalars in the fundamental representation, and gives us additional confidence in the considerations above.

In these models, SUSY is broken by

$$\langle F_{\bar{U}} \rangle \sim \left\langle \frac{\partial W_{\rm dyn}}{\partial \bar{U}} \right\rangle \sim \frac{\Lambda_{\rm comp}^2}{4\pi} (\lambda \sqrt{N})^{r-1} \left(\frac{\Lambda_{\rm lift}}{\Lambda_{\rm comp}} \right)^{3-r}. \tag{8}$$

Since $r < 3$ (otherwise the dynamical superpotential Eq. (2) does not have a good limit $\Lambda_{\rm lift} \to 0$ when $G_{\rm lift}$ is asymptotically free), we have $\langle F_{\bar{U}} \rangle \ll \Lambda_{\rm comp}^2$. The scalar components of $\bar{U}$ get a SUSY-breaking mass of order

$$m_{\bar{U}}^2 \sim \left\langle \frac{\partial^2}{\partial \bar{U}^2} \left| \frac{\partial W_{\rm dyn}}{\partial \bar{U}} \right|^2 \right\rangle \sim \frac{F_{\bar{U}}^2}{\langle \bar{U} \rangle^2} \equiv m_{\rm comp}^2. \tag{9}$$

The 'preon' fields P charged under $G_{\rm comp}$ get SUSY-breaking masses of order $m_{\rm comp}$ from effects such as

$$\Gamma_{\rm 1PI} \sim \int d^2\theta d^2\bar\theta\, \frac{16\pi^2}{\Lambda^2_{\rm comp}} \bar U^\dagger \bar U P^\dagger P \sim m^2_{\rm comp} P^\dagger P. \tag{10}$$

The scalar mass squared terms of the preons are not calculable, and therefore can have either sign. To build realistic models, we will have to assume that some of these mass-squared terms are positive. In the models we discuss, some of the fermion components of P can remain massless (in the absence of a Higgs VEV) because of unbroken chiral symmetries, and these will be identified with quarks and leptons.

With this discussion of the mass scales, we have enough information to analyze the main features of the phenomenology of these models. Masses for SM gauginos and elementary charged scalars are generated by gauge mediation from the composite scalars, so that

$$m_{\lambda,\rm SM} \sim N \frac{g^2_{\rm SM}}{16\pi^2} m_{\rm comp}, \qquad m^2_{\phi,\rm elem} \sim N \left(\frac{g^2_{\rm SM}}{16\pi^2} m_{\rm comp} \right)^2. \tag{11}$$

Note that the multiplicity factor N enhances gaugino masses compared to elementary scalar masses.

In the models we have constructed, some or all of the quarks and leptons from the first two generations are composite, while the third generation is elementary. The reason for this is that in our models the Yukawa couplings for composite quarks and leptons arise from higher-dimension operators in the fundamental theory, and are naturally small compared to one. It is difficult to accommodate the order-one top Yukawa coupling in this framework unless the top quark is elementary. Another reason for the third generation to be elementary is that stop masses of order $m_{\rm comp} \sim$ 1–10 TeV (needed to get sufficiently large gaugino masses) necessitate a large amount of fine-tuning in electroweak symmetry breaking. In order to obtain a third generation scalar mass $m_3 \gtrsim 100$ GeV we therefore require

$$m_{\rm comp} \gtrsim \frac{10 \rm TeV}{\sqrt{N}}. \tag{12}$$

We see that single sector models naturally have a superpartner spectrum similar to the 'more minimal' framework [8]. In models of this kind, there is a dangerous negative contribution to the third-generation squark masses from the heavy scalars [9], given by

$$\mu \frac{dm^2_3}{d\mu} = \frac{8g^2}{16\pi^2} C_2 \left[\frac{3g^2}{16\pi^2} m^2_{\rm comp} - m^2_\lambda \right], \tag{13}$$

where we have assumed that a single gauge group dominates and specialized to the case of two full composite generations. One way to avoid this problem is to have the compositeness scale close to 10 TeV, so that the negative contribution above does not dominate. If the compositeness scale is high, one can avoid problems if the gaugino contribution is important. From Eq. (13), we see that $m_\lambda \gtrsim m_{\rm comp}/10$ is sufficient. This requires $N \gtrsim 10$.

Most of the models we have constructed have of order N 'preonic' generations above the compositeness scale, and for $N \gtrsim 10$ the SM gauge groups are far from asymptotically free. This is compatible with perturbative unification if the compositeness scale is above (or near) the GUT scale 10^{16} GeV.

B. 'Meson' Models

We will now discuss 'meson' models where all quarks and leptons of the first two generations correspond to dimension-2 operators $P\bar U$ in the fundamental theory. This means that small Yukawa couplings involving the first two generations can be generated by:

$$W_{\rm Yuk} = \frac{1}{M^2} H (P\bar U)(P\bar U) + \frac{1}{M} H \Phi_3 (P\bar U). \tag{14}$$

where M is the scale of new physics where flavor symmetries are broken, H is a Higgs field, and Φ_3 is an elementary third-generation quark or lepton field. This gives a Yukawa matrix of the form

$$y \sim \begin{pmatrix} \epsilon^2 & \epsilon^2 & \epsilon \\ \epsilon^2 & \epsilon^2 & \epsilon \\ \epsilon & \epsilon & 1 \end{pmatrix}, \qquad \epsilon \sim \frac{\langle \bar{U} \rangle}{M}, \tag{15}$$

Additional structure is clearly needed to construct fully realistic Yukawa matrices, but for ϵ in the range 10^{-1}–10^{-2} this is a good starting point.

We will make the conservative assumption that the new physics at the scale M does not have any approximate flavor symmetries that can suppress FCNC's. In particular, this means that the Yukawa couplings λ in Eq. (1) do not conserve flavor. It is highly non-trivial that the strong dynamics in this theory nevertheless gives rise to an approximate flavor symmetry at low energies that enforces the near degeneracy of the composite scalars. The underlying reason for this is the fact that all of the composites $(P\bar{U})$ are part of a single multiplet from the point of view of the strong interactions.

Let us first consider the λ-dependent effects. The superpotential Eq. (2) depends on λ only through $\det(\lambda)$, which is flavor independent. There is nontrivial λ dependence in the effective Kähler potential, but it is proportional to $\lambda^2/(16\pi^2) \lesssim 10^{-2}$. We now consider the effects of general higher-dimension operators suppressed by the flavor scale M. The largest effects come from terms in the effective Lagrangian of the form

$$\Delta\mathcal{L}_{\text{eff}} \sim \int d^2\theta d^2\bar{\theta}\, \frac{1}{M^2} (P\bar{U})^\dagger (P\bar{U}), \tag{16}$$

which give rise to mixing between the composite generations. This translates to mixing masses between the composite generations of order

$$\frac{\Delta m^2_{jk}}{m^2_{\text{comp}}} \sim \left(\frac{\langle \bar{U} \rangle}{M} \right)^2 \sim y_{jk}. \tag{17}$$

The most stringent bounds on squark mixing come from K^0–$\bar{K}^0$ mixing, and can be summarized as

$$\text{Re}\left(\frac{\Delta m^2_{\tilde{d}\tilde{s}}}{m^2_{\text{comp}}} \right) \lesssim 10^{-1} \frac{m_{\text{comp}}}{10\text{TeV}}, \qquad \text{Im}\left(\frac{\Delta m^2_{\tilde{d}\tilde{s}}}{m^2_{\text{comp}}} \right) \lesssim 10^{-2} \frac{m_{\text{comp}}}{10\text{TeV}}. \tag{18}$$

Since $y_{ds} \sim 3 \times 10^{-4}$, this is easily satisfied even if we assume that CP violation in the flavor sector is maximal.

A striking signature of these models is that *all* scalars of the first two generations unify at the scale Λ_{comp} (which need not be close to the GUT scale). The unification holds up to effects suppressed by a loop factor, and is therefore expected to hold to 1%. This striking pattern is difficult to obtain naturally in other SUSY breaking models.

An explicit model of this type was constructed in [5] by identifying $G_{\text{lift}} \times G_{\text{comp}} \times G_{\text{global}}$ with $SU(13)_{\text{lift}} \times SU(15)_{\text{comp}} \times SU(15)_{\text{gl}}$ and taking the preon field P to consist of 3 $\Box$'s and one $\bar{\Box}$ of $SU(15)_{\text{comp}}$, where one $\Box$ of $SU(15)_{\text{gl}}$ decomposes into one SM generation.

C. 'Dimensional Hierarchy' Models

We next discuss 'dimensional hierarchy' models that explain the observed fermion mass hierarchy in terms of a hierarchy of dimensions of operators. Specifically, we assume that the first-generation quarks and leptons correspond to dimension 3 operators of the form $(P\bar{U}\bar{U})$, second-generation quarks and leptons correspond to dimension 2 operators $(P\bar{U})$, and third generation quarks and leptons are elementary (dimension 1). In this case, Yukawa couplings involving the composite states arise from terms in the tree-level superpotential of the form

$$W_{\text{Yuk}} = \tfrac{1}{M^4} H(P\bar{U}\bar{U})(P\bar{U}\bar{U}) + \tfrac{1}{M^3} H(P\bar{U}\bar{U})(P\bar{U}) + \tfrac{1}{M^2} H\Phi_3(P\bar{U}\bar{U}) \tag{19}$$

$$+ \tfrac{1}{M^2} H(P\bar{U})(P\bar{U}) + \tfrac{1}{M} H\Phi_3(P\bar{U}), \tag{20}$$

giving rise to a Yukawa matrix of the form

$$y \sim \begin{pmatrix} \epsilon^4 & \epsilon^3 & \epsilon^2 \\ \epsilon^3 & \epsilon^2 & \epsilon \\ \epsilon^2 & \epsilon & 1 \end{pmatrix}, \qquad \epsilon \sim \frac{\langle \bar{U} \rangle}{M}. \tag{21}$$

This structure reproduces the main features of the observed fermion mass hierarchy for $\epsilon \sim 10^{-1}$. A striking signature of these models is that the first- and second-generation scalars unify in two multiplets at the scale Λ_{comp}.

In this scenario there is no approximate flavor symmetry at low energies because the first- and second-generation fields belong to different strong-interaction multiplets. We therefore have

$$\frac{\Delta m^2_{\tilde{d}\tilde{s}}}{m^2_{\text{comp}}} \sim \sin\theta_{\text{c}} \sim 10^{-1}. \tag{22}$$

Comparing with the bounds from the K^0–$\bar{K}^0$ system Eq. (18), we see that $m_{\text{comp}} \sim 10$ TeV is sufficient to suppress FCNC's but we require either a 10% fine-tuning or a 10% suppression of *CP*-violating effects in the squark masses.

An explicit model of this type was constructed in [5] by identifying $G_{\text{lift}} \times G_{\text{comp}} \times G_{\text{global}}$ with $SU(2)_{\text{lift}} \times SU(16)_{\text{comp}} \times SU(16)$ and taking the preon field P to consist of 2 $\square$'s and one antisymmetric tensor of $SU(16)_{\text{comp}}$.

D. A Simple but Speculative Model

We now present a simple model whose dynamics we do not know how to analyze completely. If we make a reasonable dynamical assumption, this model gives rise to compositeness and SUSY breaking by the mechanism discussed in above. The particle content is:

	$SU(k)$	$SO(10)$	$SU(10)$	$SU(2)$
Q	$\square$	$\square$	**1**	**1**
L	$\bar{\square}$	**1**	$\square$	**1**
$\bar{U}$	**1**	$\square$	$\bar{\square}$	**1**
P	**1**	**16**	**1**	$\square$

.

For $\langle \bar{U} \rangle \gg \Lambda_{10}$, $SO(10) \times SU(10)$ is broken to the diagonal $SO(10)$ subgroup and $SU(k)$ gaugino condensation gives rise to a dynamical superpotential

$$W_{\text{dyn}} \sim \bar{U}^{10/k}. \tag{23}$$

The potential therefore slopes toward $\bar{U} \to 0$ for $k < 10$.

The dynamics for small values of $\langle \bar{U} \rangle$ involves the strong-coupling behavior of the $SO(10)$ gauge theory with spinors, which is presently not well understood. The $SO(10)$ gauge theory has a dual description in terms of an $SU(2) \times SU(7+k)$ gauge theory [10]; this dual is not weakly coupled in the infrared, so we cannot use it to determine the behavior of the Kähler potential for $\bar{U}$. Based on analogies with similar duals, one expects this theory to have an infrared fixed point [10,11].

If we assume that the anomalous dimension of $\bar{U}$ is sufficiently large (> 0.1), then there is a local SUSY-breaking minimum. The fermions from the **16**'s are exactly massless far from the origin, and because there can be no phase transitions as a function of moduli they are massless at the local minimum as well. This model therefore contains two composite fermionic **16**'s which can be identified with two SM generations (with right-handed neutrinos) if we embed the SM into $SU(10)$ via the standard $SO(10)$ GUT embedding FCNC's are suppressed by the approximate global $SU(2)$ symmetry of the strong dynamics. Above the compositeness scale, this model has $3 + k/2$ additional 'preonic' generations.

Yukawa couplings for the composite generations can be induced if we include a Higgs field, H embedded in the $\square$ of the global $SU(10)$ by operators of the form

$$W_{\text{Yuk}} = \frac{1}{M} P P H \bar{U} \tag{24}$$

This gives Yukawa couplings $y \sim \langle \bar{U} \rangle / M$ for the composite quarks and leptons. (Compared with our previous expressions, this corresponds to the composite operators being dimension $\frac{3}{2}$.) Thus the flavor scale M can be pushed up even higher in this model, and FCNC's are even more suppressed than in our 'meson' models.

III. CONCLUSIONS

We have seen that there is a wide class of realistic models which dynamically break SUSY and produce composite quarks and leptons, all in a single (strongly-coupled) sector. These models are remarkably simple; many of the fermion mass hierarchies follow naturally, and approximate flavor symmetries of the strong interactions can guarantee the natural suppression of flavor-changing neutral currents (including ϵ_K) with no fine-tuning.

All of these models have a very distinctive phenomenology: the composite sfermions of the first two generation are heavier than the gauginos, Higgsinos, and third-generation sfermions; and the composite sfermions unify at the compositeness scale.

A (perhaps) surprising result of our recent work [5] is the wide variety of dynamics that can give rise to simultaneous compositeness and SUSY breaking. We have found models whose low-energy dynamics is governed by either confinement, a non-trivial infrared fixed point, or a free-magnetic phase. We believe that it is quite likely that further progress in understanding the dynamics of SUSY gauge theories will lead to the discovery of many additional models that display the dynamics illustrated here.

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