Dynamical supersymmetry breaking at low energies

Michael Dine
Santa Cruz Institute for Particle Physics, University of California, Santa Cruz, California 95064

Ann E. Nelson
Department of Physics, 9500 Gilman Drive 0319, University of California, San Diego, La Jolla, California 92093-0319
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Conventional approaches to supersymmetric model building suffer from several naturalness problems: they do not explain the large hierarchy between the weak scale and the Planck mass, and they require fine-tuning to avoid large flavor-changing neutral currents and particle electric dipole moments. The existence of models with dynamical supersymmetry breaking, which can explain the hierarchy, has been known for some time, but efforts to build such models have suffered from unwanted axions and difficulties with asymptotic freedom. In this paper we describe an approach to model building with supersymmetry broken at comparatively low energies which solves these problems, and give a realistic example.

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I. INTRODUCTION

In recent years, supersymmetry has emerged as a leading candidate for the origin of electroweak symmetry breaking. Supersymmetry offers a cure to the hierarchy problem (or at least to the problem of quadratic divergences) and, in a desert scenario, yields a successful unification of coupling constants. Extensive effort has been devoted to the phenomenology of what has become known as the "minimal supersymmetric standard model" (MSSM). This model has no difficulty accommodating present experimental constraints. The features of the MSSM are often understood, or at least motivated, by considering $N=1$ supergravity theories.

Yet there are a number of reasons to be concerned about this rosy picture; these concerns can also be motivated within the framework of supergravity theories. First, the supersymmetry-breaking scale is put in by hand. The promise of supersymmetry to *explain* the hierarchy is unfulfilled. Furthermore, the masses of supersymmetric partners are not calculable, and must also be put in by hand to be above the experimental bounds. Moreover, in these $N=1$ models, the new physics associated with the "hidden sector" responsible for supersymmetry breaking is completely inaccessible. This last objection is not fundamental, but it is a disappointing feature of these theories. Finally, and perhaps most seriously, there is the "flavor problem." The pattern of soft breakings in these theories is highly restricted by flavor-changing neutral currents. Usually it is simply assumed that at some high-energy scale, the squarks and sleptons are degenerate. In the framework of standard supergravity theories, however, there is no reason for such relations to hold, and indeed they do not in generic superstring compactifications [1]. More generally, it is hard to understand how theories in which the quark and squark masses are generated at some very high-energy scale can give rise to significant squark and slepton degeneracy [2]. There have been many speculations on solutions to all of these problems; still the picture is not completely satisfying theoretically, and most of these speculations involve physics which is experimentally out of reach for the foreseeable future.

In this paper we wish to consider an alternative picture of supersymmetry breaking, which has not been considered since the earliest days of supersymmetry model building. We wish to explore the possibility that supersymmetry is dynamically broken, by new physics associated with (multi) TeV energies. We will construct such a model, where the squark and slepton masses are calculable. The model has many desirable features; in fact, it will solve all the problems listed above. It will also make some predictions, not only about the size of soft breaking terms, but also about the particle content at the weak scale. In particular, beyond the particles of the MSSM, it predicts the existence of at least one singlet and a set of mirror quarks and leptons, all at experimentally accessible energies.

Models exhibiting dynamical supersymmetry breaking (DSB) have been known for some time [3]. The authors of Ref. [4] attempted to construct models with supersymmetry broken in the multi-TeV energy range, but ran into a variety of problems. Two problems, in particular, seemed generic: there were light Goldstone bosons, or axions, and QCD was not asymptotically free, hitting its Landau singularity a few decades above the scale of the new, supersymmetry-breaking physics. In the present work, we will exhibit a model which solves both problems. The would-be axion will gain mass as a result of another strong group in addition to QCD, and the model will be structured so that QCD is nearly asymptotically flat. Ordinary particles, quarks, leptons, gauge bosons, and their superpartners, "learn of" supersymmetry breaking through gauge interactions. As a result, there is automatically sufficient degeneracy in the squark and slepton spectrum to insure adequate suppression of

flavor-changing neutral currents. $SU(2)\times U(1)$ breaking will arise through loop corrections to Higgs boson masses involving top quarks, in a manner discussed long ago [5]. Breaking the electroweak symmetry without fine-tuning will also require the presence of a light $SU(2)\times U(1)$ singlet, and new superfields with vectorlike interactions.

The model we will present here is meant as an existence proof. It has certain drawbacks. None of these are fatal, nor is it clear that any of them are generic. Indeed, we strongly suspect that a more elegant model is lurking somewhere. Perhaps the most serious problem is just that the model is rather complicated, involving four additional gauge groups. It is not, however, nearly as complicated as recent proposals for technicolor models, and the symmetry group is not larger than some encountered in string compactifications. Moreover, unlike technicolor models, it is not necessary for large numbers of groups to become strong within a few decades of one another. Apart from the group actually responsible for breaking supersymmetry, there is one other group which cannot be too weak; otherwise there is a light Goldstone boson which is inconsistent with the red giant and supernova limits. We will see that at high energies this requires at most a very mild fine-tuning. There is also the potential for generating a large Fayet-Iliopoulos D term for hypercharge. Solving this will require an approximate equality of certain gauge couplings. This equality will be seen to be "natural" in the sense that it does not receive large radiative corrections. Such an equality could arise within the framework of grand unification or superstring theory. Finally, there are potential cosmological problems which could be solved by higher-dimension operators: domain walls and long-lived massive states. Still, the model has virtues: a hierarchy between the weak scale and the shortest distance scales is naturally generated, flavor-changing processes and new sources of CP violation are naturally suppressed, and it is otherwise consistent with all present day experiments. Of course, the cosmological constant problem remains as a most troublesome naturalness issue.

II. DYNAMICAL SUPERSYMMETRY BREAKING

The hierarchy problem has two aspects, both of which one might hope to address within the framework of supersymmetry. One is the problem of quadratic divergences of scalar masses. The second is the existence of an extremely small dimensionless number, which we can think of as the ratio of the weak scale to M_P or some unification scale. It was Witten who first clearly stated in what sense supersymmetry might solve this second problem [6]. He noted that (in the case of global supersymmetry) any small vacuum energy signifies supersymmetry breaking. Yet if supersymmetry is unbroken at the classical level, it remains unbroken to all orders in perturbation theory as a consequence of nonrenormalization theorems. However, the proof of this statement is inherently perturbative, and the result need not hold beyond perturbation theory. Thus effects of order e^{-a/g^2}, where g is some coupling constant, might give rise to supersymmetry breaking and explain the large hierarchy. Witten also formulated a set of conditions under which supersymmetry breaking might or might not occur. Most important of these was the existence (in perturbation theory) of a massless fermion which could play the role of a Goldstone fermion. He also showed that a certain index (the "Witten index") must vanish if supersymmetry is to be broken, and showed this index to be nonzero in a number of interesting cases [7].

Subsequent work showed that in many cases, nonperturbative effects do violate the nonrenormalization theorems. In some cases these are due to instantons and can be calculated explicitly in a systematic semiclassical expansion; in some cases they can be understood in terms of other nonperturbative effects, such as gluino condensation. It will be helpful for what follows to review the results for "supersymmetric QCD" [8]. For our purposes, this is a theory with gauge group $SU(N)$ with N_f chiral multiplets in the N representation, Q_f, and N_f in the $\bar{N}$ representation, $\bar{Q}_f$. Before including a mass term or other superpotential term for the "quark" fields, the theory has a nonanomalous $SU(N_f)\times SU(N_f)\times U(1)_V\times U(1)_R$ symmetry, where the last two symmetries are a baryon-number-like transformation and an R transformation with charges chosen to avoid anomalies. This model also contains, at the classical level, a large set of degenerate vacuum states, described by several parameters, referred to as "flat directions" of the potential. These are just directions in which the D terms vanish. It is not hard to convince oneself that up to symmetry transformations, for $N_f<N$ the most general zero-energy state is described by the expectation values:

$$Q=\begin{pmatrix} a_1 & 0 & \cdots \\ 0 & a_2 & \cdots \\ \cdots & 0 & a_{N_f} \\ 0 & \cdots & 0 \\ \cdots & \cdots & \cdots \\ 0 & \cdots & 0 \end{pmatrix}=\bar{Q}\ . \tag{2.1}$$

The case where $N_f=N-1$ is the easiest to analyze. In this case, in these flat directions, the gauge symmetry is completely broken. Moreover, by choosing the vacuum expectation values (VEV's) large enough, one can make the theory arbitrarily weakly coupled [the effective coupling is $\alpha(M_V)$, the asymptotically free coupling of the theory evaluated at the scale of the gauge boson masses]. In these vacua the light degrees of freedom are a set of Goldstone bosons and their superpartners. These can be described by a matrix-valued field $\Phi_{f,\bar{f}}=Q_f\bar{Q}_{\bar{f}}$. In these vacua, in order to determine if supersymmetry breaking occurs, one must compute the effective action for these light degrees of freedom. In particular, supersymmetry breaking requires that one generate a superpotential for these fields. Any superpotential must respect the original flavor symmetries of the theory. The $SU(N_f)\times SU(N_f)$ symmetry implies that the action must be a function only of $\det(\Phi)$. The $U(1)_R$ symmetry then determines the form of the superpotential uniquely:

$$W=\frac{\Lambda^{2N+1}}{\det(\Phi)}\ , \tag{2.2}$$

where Λ is the scale parameter of the SU(N) group. A completely straightforward instanton calculation precisely yields the various component field interactions implied by this Lagrangian [8].

For $N_f<N-1$, one can repeat most of the analysis above. In particular, the symmetries determine the form of any superpotential uniquely, in terms of the $N_f\times N_f$ field, $\Phi_{f\bar{f}}=Q_f\bar{Q}_{\bar{f}}$:

$$W=b\Lambda^{(3N-N_f)/(N-N_f)}\det(\Phi)^{-1/(N-N_f)}\ . \tag{2.3}$$

A somewhat different analysis is required to show that this superpotential is generated in this case. For $N_f>N-1$, no superpotential is generated.

What are the implications of this superpotential? These examples illustrate that the nonrenormalization theorems *do* break down nonperturbatively. However, at least at weak coupling, these theories do not have a good ground state, and at best admit a cosmological interpretation. If one adds mass terms for the quarks, the potential is stabilized, and one finds N supersymmetric ground states, in agreement with Witten's calculation of the index. If one wants to find theories in which supersymmetry is broken and one has a good ground state, it is necessary to study chiral theories. In Ref. [4], a number of chiral theories were studied which do exhibit DSB. The main conditions for this are (1) the absence of flat directions in the classical theory and (2) the existence of a nonanomalous, continuous symmetry which is spontaneously broken. These conditions are not hard to understand. The second implies the existence of a Goldstone boson. Unbroken supersymmetry would imply the existence of a *scalar* partner for this field, but this would imply the existence of a flat direction, contradicting (1). One can imagine loopholes to this argument, but these conditions seem to be a good guide to finding theories with DSB.

III. STRATEGIES FOR MODEL BUILDING

Having established the existence of models with dynamical supersymmetry breaking, it is natural to try and build realistic models of low-energy supersymmetry incorporating it. There are two strategies one might adopt. First, one might use these models as hidden sectors for $N=1$ supergravity models. Aspects of this problem have been discussed elsewhere [9]. However, even if this program is successful, it has little predictive power; the superparticle spectrum remains a function of unknown parameters, and the origin of degeneracy among squarks and sleptons remains mysterious.

Alternatively, one can consider the possibility that supersymmetry is broken at comparatively low energies, and that the breaking of supersymmetry is fed to the superpartners of ordinary fields through gauge interactions. The most straightforward (though perhaps not the most clever) way to proceed is to take a model of the type discussed above with DSB, and gauge a global symmetry, identifying it with one of the usual gauge interactions. The squarks, sleptons, and gauginos will gain mass through loop effects. Previous efforts to realize this scenario floundered on two problems. First, there is often a problem with the asymptotic freedom of QCD. Models with DSB and large enough global symmetry groups to gauge an SU(3) subgroup typically contain large numbers of triplets and antitriplets. For example, in Ref. [4], the simplest such model had gauge group SU(11) and gave rise to 11 new flavors of quarks, while in a supersymmetric theory if one requires that QCD not have a Landau pole below the unification scale of 10^{16} GeV at most four new flavors of quarks are allowed at the weak scale.

A second problem is the existence of unacceptable Goldstone bosons or axions. As we have noted, all known examples of dynamical supersymmetry breaking require the presence of a spontaneously broken global symmetry, and with it a massless Goldstone boson. In many cases, this boson is an axion, once QCD is taken into account. In any case, the decay constant of this boson is of multi-TeV order, and thus it is typically inconsistent with astrophysical limits.

Here we will describe a model which avoids both of these problems. The dangerous Goldstone boson will gain mass as a result of anomalies with respect to another strong gauge interaction, known as "R color." The problem of nonasymptotic freedom will be avoided by using a slightly more complicated strategy than that described above. The extra gauged symmetry in the "supercolor sector" (i.e., the sector responsible for breaking supersymmetry) will not be identified with the standard model gauge interactions, but with R color. There will be some additional fields carrying R color as well as ordinary gauge quantum numbers, which will cancel R-color anomalies and act as the "messengers" of supersymmetry breaking. As a result, QCD will be only barely nonasymptotically free.

One unpleasant feature of the model, which differs from earlier models of DSB, is that it has classically flat directions. The degeneracy is lifted by nonperturbative effects, but a supersymmetric minimum appears at infinite value of some scalar fields; the theory has no ground state. However we can find a *local* minimum of the potential which violates supersymmetry. For small coupling, this minimum will be essentially stable against tunneling. We will not worry here about how the Universe might have found itself in this state.

IV. THE MODEL

A. Fields and Lagrangian

Let us now turn to the actual model. Apart from SU(3)×SU(2)×U(1), the gauge group of the model is

$$\mathrm{SU}(7)\times\mathrm{SU}(2)\times\mathrm{SU}(3)_L\times\mathrm{SU}(3)_R\ . \tag{4.1}$$

The SU(7)×SU(2) groups will be referred to as "supercolor." The SU(7) gets strong, and nonperturbative terms in the superpotential generated by the SU(7), in conjunction with the D terms from the SU(2), will be responsible for supersymmetry breaking. The matter fields of the

model consist, first, of the usual quark and lepton superfields, a pair of Higgs boson doublets, and a singlet S. The latter particle will be necessary for achieving $SU(2)\times U(1)$ breaking. To describe the additional fields of the model, it is convenient, first, to ignore the standard model fields and interactions, and impose a global SU(7) symmetry. The usual gauge interactions will lie within this SU(7). This procedure will allow us to turn immediately to the essential dynamical features of the model. Later we will return to the realistic situation where the global SU(7) symmetry is explicitly broken to an $SU(3)\times SU(2)\times U(1)$ gauged subgroup. Under

$$SU(7)\times SU(2)\times SU(3)_L\times SU(3)_R\times SU(7)_G ,$$

the additional fields are

$$\begin{aligned} &Q=(7,1,\bar{3},1,1),\quad \bar{Q}=(\bar{7},1,1,3,1) ,\\ &q=(7,2,1,1,1),\quad \bar{u}=(\bar{7},1,1,1,1) ,\\ &\bar{d}=(\bar{7},1,1,1,1) ,\\ &X=(1,1,3,\bar{3},1),\quad \bar{X}(1,1,\bar{3},3,1) ,\\ &f=(1,1,3,1,7),\quad \bar{f}=(1,1,1,\bar{3},\bar{7}) ,\\ &l=(1,2,1,1,1) . \end{aligned} \tag{4.2}$$

The superpotential of the model is

$$\begin{aligned} W=&\lambda_1\bar{Q}XQ+(\lambda_2/3)\det(X^3)+\lambda_3 q\bar{u}l\\ &+(\lambda_4/3)\det(\bar{X}^3)+\lambda_5\bar{f}\bar{X}f . \end{aligned} \tag{4.3}$$

Note that the model is anomaly-free. When we return to consider the ordinary $SU(3)\times SU(2)\times U(1)$, we will simply embed this group in the standard way in an SU(5) subgroup of SU(7). We will thus take f and $\bar{f}$ each to breakup into a triplet, a doublet, and two singlets. The coupling λ_5 will then actually represent four independent parameters, which we refer to as $\lambda_5^{t,d,s,s'}$.

We will suppose that SU(7) is the strongest group (i.e., the one with the largest Λ parameter), and that all of the couplings in the superpotential are small. In this approximation, the SU(7) sector of the theory is an example of supersymmetric QCD with seven colors and five flavors. Grouping the 7 and $\bar{7}$'s of the theory into fields Q and $\bar{Q}$, this theory has flat directions of the form of Eq. (2.1) [it is necessary to use the approximate $SU(5)\times SU(5)$ flavor symmetry to bring these fields to this form]. Nonperturbatively as described above, a superpotential of the form

$$W_{\rm np}=\Lambda^8/(\det\Phi)^{1/2} \tag{4.4}$$

is generated, where

$$\Phi=(\bar{Q}_{\bar{f}}Q_f) . \tag{4.5}$$

For small couplings λ_i, we want to study the potential

$$W=W_{\rm cl}+W_{\rm np} \tag{4.6}$$

(and the D terms for the various groups). We expect that for small λ_i, the minimum of the potential will lie at large values of the fields, justifying this analysis. Provided $\lambda_i \ll g_a$ (the gauge couplings), we should be able to find the minima by looking at flat directions of the D terms. Of course, the full theory does not have the $SU(5)\times SU(5)$ symmetry used above, and so one must consider a more general set of flat directions of the D terms. The structure of the complete potential is quite complicated, and we will not be able to survey the entire field space. Instead, we will look for a local minimum of the potential in a particular direction in the field space. This will be described by the expectation values

$$\begin{aligned} &Q=\begin{pmatrix} a & 0 & 0\\ 0 & a & 0\\ 0 & 0 & a\\ 0 & \cdots & 0\\ \cdots & \cdots & \cdots \end{pmatrix}=\bar{Q},\quad X=\mathrm{diag}(x,x,x)\\ &q=\begin{pmatrix} 0 & 0\\ 0 & 0\\ 0 & 0\\ b & 0\\ 0 & c\\ 0 & 0\\ 0 & 0 \end{pmatrix},\quad \bar{u}=\begin{pmatrix} 0\\ 0\\ 0\\ d\\ e\\ 0\\ 0 \end{pmatrix},\quad \bar{d}=\begin{pmatrix} 0\\ 0\\ 0\\ f\\ g\\ 0\\ 0 \end{pmatrix},\quad L=\begin{pmatrix} h\\ 0 \end{pmatrix} \end{aligned} \tag{4.7}$$

with

$$\begin{aligned} &|f|^2=|b|^2-|d|^2,\quad |g|^2=|c|^2-|e|^2 ,\\ &|h|^2=|c|^2-|b|^2 . \end{aligned} \tag{4.8}$$

We establish that there exist a *local* minimum of the potential of this form. Note that this minimum leaves over an SU(3) gauge symmetry which is a linear combination of the $SU(3)_L$, $SU(3)_R$, and an SU(3) subgroup of SU(7). Thus SU(3), known as R color, will subsequently also get strong. There is an unbroken SU(2) subgroup of the SU(7), which will play no role in the subsequent discussion. The fields f, $\bar{f}$, and $\bar{X}$ play no role in the dynamics which break supersymmetry; but simply serve as "messenger" particles which communicate supersymmetry breaking to ordinary superfields.

The rest of the model, which we will refer to as the ordinary sector, just consists of the MSSM, without the bilinear μH_1H_2 term in the superpotential (this can be eliminated by imposing a Z_3 discrete symmetry). We get rid of this term because otherwise electroweak symmetry breaking would require an unacceptable fine-tuning of μ. Instead, breaking $SU(2)\times U(1)$ will require the addition of a gauge singlet S and vector like superfields $D,\bar{D},L,\bar{L}$ transforming as $(3,1,-\frac{1}{3})(\bar{3},1,\frac{1}{3}),(2,1,\frac{1}{2}),(2,1,-\frac{1}{2})$ under the ordinary $SU(3)\times SU(2)\times U(1)$ interactions.

B. Overview of symmetry-breaking in the model

Before going through the detailed analysis of the model, we summarize the basic features. Having established that the minimum of the potential is of the form of Eq. (4.7), one of our main goals is to determine the masses of squarks, sleptons, and gauginos, as well as Higgs boson particles. These will arise as a result of the gauge couplings of the f and $\bar{f}$ fields. The scalar components of

these fields, as well as the fields $\bar{X}$, can gain mass at one loop through graphs such as those shown in Fig. 1. It turns out, that for a range of parameters, the mass squared of these fields are negative. Minimizing the resulting potential yields a VEV for $\bar{X}$ of the form

$$\bar{X}=\mathrm{diag}(\bar{x},\bar{x},\bar{x}),\quad f=\bar{f}=0\ . \tag{4.9}$$

The fields f and $\bar{f}$ receive an additional contribution to their mass from the VEV of $\bar{X}$. Below the scale of these VEV's, one has an unbroken SU(3) gauge theory without matter fields at all [the S(3) gauginos gain mass at one loop].

Ordinary gauginos gain mass through diagrams of the type shown in Fig. 2. Squarks and sleptons gain masses through diagrams of the type shown in Fig. 3. For a range of parameters, these contributions can be shown to be positive. Note that the masses of squarks, sleptons, and gauginos depend in a simple way on their gauge couplings. Squarks are generically heaviest, lepton doublets and Higgs bosons are lighter by roughly a factor α_2/α_3, and singlet leptons are lightest. In order that $\mathrm{SU}(2)_L$ be broken, it is necessary that one Higgs particle obtain a negative mass squared. This can occur for the Higgs boson which couples to top quarks [5], as a result of the diagram of Fig. 4. While this diagram is nominally one higher order in the loop expansion, it is enhanced by the fact that the squark masses are larger than the doublet masses by a factor α_3/α_2, and by a logarithm, and can be larger than the positive two-loop contributions. For a range of parameters, as a result, the Higgs boson mass squared can be negative. In this model, however, there is no H_1H_2 term in the potential, so in order that there be a suitable quartic coupling for the Higgs boson field, it is necessary to include a singlet field, with couplings SH_1H_2 and S^3. Furthermore, in order to obtain a sensible breaking of $\mathrm{SU}(2)\times\mathrm{U}(1)$ with masses for all the quarks and leptons and without fine tuning it will be necessary to add new vectorlike fermions carrying $\mathrm{SU}(3)\times\mathrm{SU}(2)\times\mathrm{U}(1)$ gauge charges, which couple to the singlet and gain mass from its VEV.

Finally, we have to worry about the various global symmetries of the model. The vectorlike symmetries are preserved at this minimum, but the $\mathrm{U}(1)_R$ symmetry is broken, giving rise to an axion. This symmetry has no SU(7) anomaly. However, the axion does get a mass from the unbroken SU(3) R color, which is of the order of the scale at which this group gets strong, squared, and divided by its decay constant. This mass can easily be of the order 10 MeV, so its production in stars can be adequately suppressed. There remain a number of cosmological worries about this model; these include domain walls and stable massive particles, and will be dealt with later, as will the question of Fayet-Iliopoulos terms. We now turn to a detailed discussion of each of these points.

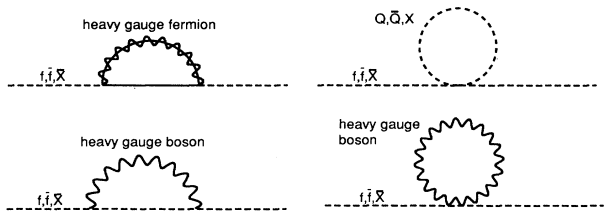

FIG. 1. One-loop diagrams contributing to the masses of the $\bar{X}$, f, and $\bar{f}$ fields.

FIG. 2. One-loop diagrams contributing to the masses of the gauginos.

C. Supersymmetry breaking

First, let us turn to the problem of minimizing the potential. It is not hard to see that the minimum in the direction of Eq. (4.7) cannot be the global minimum; the global minimum has zero energy. At the classical level the theory has a flat direction with

$$X=\mathrm{diag}(x,0,0)=\bar{X}\ , \tag{4.10}$$

all other fields vanishing. Once one considers the nonperturbative piece of the superpositional, this direction is no longer flat. However, it is possible to let $X\to\infty$, $Q\to\infty$, and the fields q, $\bar{u}$, $\bar{d}$, and L tend to zero in such a way that the total energy tends to zero. Note first that for large x, the unbroken symmetry is an $\mathrm{SU}(2)\times\mathrm{SU}(2)\times\mathrm{U}(1)$. One Q flavor gains mass, as does a (2,2) field from X. After integrating out massive states, there is no dimension-four term in the superpotential for the light Q's. In order to minimize the nonperturbative contribution to the potential, then, one wants to let Q get large (though not as fast as X), while the other fields get smaller more slowly. For example, the scaling

$$Q\sim x^{1/4}\sim\bar{Q},\quad q\sim x^{-1/8}\sim\bar{u}\sim\bar{d}$$

gives an energy tending to zero as $x^{-1/2}$. As explained in the Introduction we will not worry about the global

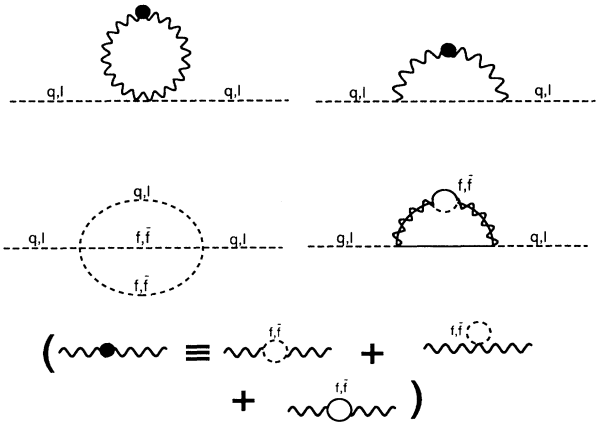

FIG. 3. Two-loop diagrams contributing to squark, slepton, and Higgs boson masses.

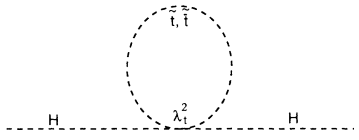

FIG. 4. Loop correction to Higgs boson mass in the low-energy effective theory, which gives negative contribution proportional to the square mass squared.

structure of the potential and simply assume that somehow the Universe finds itself in the vacuum of interest, and does not tunnel out.

Let us turn to the problem of minimizing the potential in the direction of Eq. (4.7). The problem is easiest to analyze in the limit $\lambda_3 \ll \lambda_1 \ll \lambda_2$. In the limit $\lambda_3 \to 0$, the auxiliary fields F_X and F_Q should vanish. $F_X = 0$ gives

$$x^2 = (\lambda_1/\lambda_2)a^2 \tag{4.11}$$

while $F_Q = 0$ gives an expression for a in terms of $\det\Phi$:

$$a = \tfrac{1}{2}\lambda_2^{1/2}\lambda_1^{-3/2}\Lambda^8\tilde{\Phi}^{-1/2} , \tag{4.12}$$

where we have defined

$$a^4 \det\Phi = \tilde{\Phi} . \tag{4.13}$$

Now one can plug this expression for a into the remaining terms in the potential:

$$\left|\frac{\partial W}{\partial q}\right|^2 + \left|\frac{\partial W}{\partial \bar{u}}\right|^2 + \left|\frac{\partial W}{\partial \bar{d}}\right|^2 + \left|\frac{\partial W}{\partial l}\right|^2 . \tag{4.14}$$

We can obtain the dependence of the VEV's $b - h$ on the couplings $\lambda_1, \ldots, \lambda_3$ by scaling arguments. A simple exercise gives

$$(b,c,d,e,f,g,h) \sim \lambda_3^{-1/4}\lambda_1^{3/16}\lambda_2^{-1/16}\Lambda . \tag{4.15}$$

One can check that, for a finite range of parameters, the minimum of the potential is indeed of this form, with an unbroken SU(3).

For finite but small λ_3, the relations

$$F_Q = 0, \quad F_X = 0 , \tag{4.16}$$

are not exactly satisfied. To determine the corrections, we need to compute the shifts in the VEV's a and x to the next nontrivial order in λ_3. Again, it is not hard to determine how these scale with couplings. The shift in a, δa, can be determined by computing the a tadpole, $\partial V/\partial a$, and dividing by the a mass squared. Using our scaling results above, one finds that

$$\delta a \sim \Lambda\lambda_2^{7/16}\lambda_3^{3/4}\lambda_1^{-21/16} .$$

δx is smaller by a factor $\sqrt{\lambda_1/\lambda_2}$. We can estimate F_X and F_Q by writing

$$F_X = \frac{\partial^2 W}{\partial a\,\partial x}\delta a + \frac{\partial^2 W}{\partial x^2}\delta x \tag{4.17}$$

with a similar equation for F_Q. Note that the second derivatives here are just elements of the lowest-order mass matrix. In the limit $\lambda_1 \ll \lambda_2$, one finds that F_X is largest:

$$F_X \sim \Lambda^2\lambda_2^{13/24}\lambda_3^{5/6}\lambda_1^{-5/8} \gg F_Q \sim \Lambda^2\lambda_2^{1/24}\lambda_3^{5/6}\lambda_1^{-1/8} . \tag{4.18}$$

Note that the spectrum of particles in the supercolor sector is nearly supersymmetric. The breaking of supersymmetry is represented by the small values of the F components (small by powers of the couplings in the superpotential), which give rise to small splittings within the multiplets. The gauge bosons associated with the broken $SU(3)_L \times SU(3)_R$ are also nearly supersymmetric. The spectrum is simpler to work out if the gauge couplings of these groups, g_L and g_R are identical; as we have already remarked in the Introduction, this condition must, in fact, be satisfied if the model is to be realistic. The expectation value a is larger than that of x, so, neglecting the x VEV there are two massive gauge bosons, with mass squared $g_L^2 a^2$, and $(2g_7^2 + 2g_7 g_L)a^2$, and one massless eigenstate corresponding to the unbroken SU(3). If we assume $g_L \ll g_7$, the former is the lighter state; it is simply the linear combination

$$B^\mu = (1/\sqrt{2})(A_L^\mu + A_R^\mu) . \tag{4.19}$$

This hierarchy of VEV's will be important when we estimate the loop contributions to various masses. At the classical level there are many massless states, such as f, $\bar{f}$, and $\bar{X}$. To determine whether these fields obtain expectation values, one needs to compute their masses. These will arise from the one-loop diagrams shown in Fig. 1. In the limit in which we are working, in which supersymmetry breaking is small, one can evaluate the masses perturbatively in powers of F_X. This is conveniently done using supergraph techniques. The required diagrams are then indicated in Fig. 5. Because of the hierarchy of VEV's, it is not necessary to consider diagrams such as that of Fig. 6, with external X's; it is also not necessary to consider diagrams with external F_Q's. To evaluate the diagrams it is convenient to choose the supersymmetric analog of R_ξ gauge [10]. In this gauge, the gauge propagator is simply

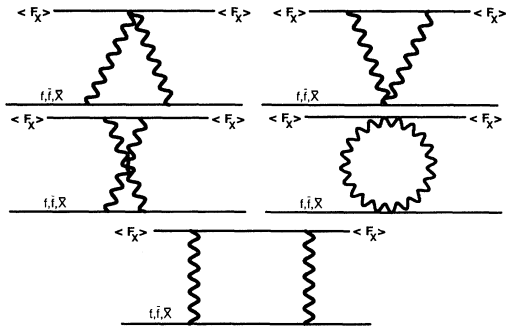

FIG. 5. Supergraphs contributing to the scalar masses. Solid lines are chiral superfields; wavy lines denote gauge fields. X's denote vacuum insertions.

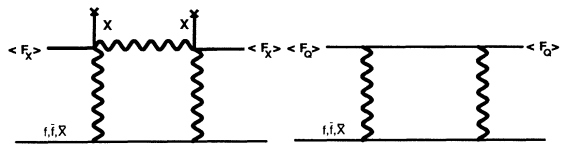

FIG. 6. Examples of diagrams suppressed by powers of couplings.

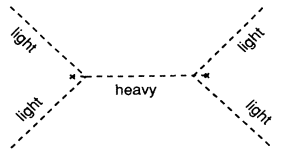

FIG. 7. In the supersymmetric limit, D terms in the potential associated with broken gauge generators are canceled by exchange of the massive scalars in the vector multiplet.

$$\Delta=\frac{\delta^4(\theta_1-\theta_2)}{p^2-M_V^2} . \tag{4.20}$$

The θ integrations are trivially performed, and one obtains for the scalar mass

$$m_S^2=-\mathcal{C}\frac{g^4}{16\pi^2}\frac{F_X F_X^\dagger}{M_V^2} . \tag{4.21}$$

$\mathcal{C}$ is a group-theory factor which is easy to work out in each case. For example, under the surviving SU(3), $\bar{X}$ decomposes as a singlet, x^s and an octet, x^a; for these, $\mathcal{C}$ equals $\frac{8}{3}$ and $\frac{7}{6}$, respectively. For the f and $\bar{f}$ fields, all of which are triplets, $\mathcal{C}=\frac{2}{3}$.

We wish to determine the pattern of symmetry breaking at this stage. In particular, we will ask if the effective potential has a local minimum at which SU(3) remains unbroken; this requires that only the singlet x^s obtain a VEV. To investigate this, we need to determine the form of the quartic terms in the potential, which arise from two sources. First, there are the terms in the original superpotential. In terms of canonically normalized fields, this superpotential takes the form

$$W=\frac{\lambda_4}{\sqrt{3}}\left[\frac{x^{s3}}{3}-\frac{1}{2}x^s x^{a2}+O(x^{a3})\right] , \tag{4.22}$$

where

$$\bar{X}=x^s/\sqrt{3}+\sqrt{2}x^a T^a . \tag{4.23}$$

If this were the end of the story, it is easy to check that the SU(3)-preserving extremum of the potential (including the loop-generated mass terms) is unstable. If one simply looks for an extremum with $x^s\neq 0$, $x^a=0$, one finds that the octet masses are tachyonic.

However, this is *not* the whole story; there are additional tree level supersymmetric-breaking quartic couplings in the effective low-energy theory which describes the $\bar{X}$ fields. To understand this, consider the terms in the potential of the full theory associated with the auxiliary D fields for $SU(3)_L$ and $SU(3)_R$. These terms are nonvanishing for $\bar{X}$ fields of the type we are describing [remember that $\bar{X}$ transforms as a $(\bar{3},3)$]. If supersymmetry were unbroken, this would be irrelevant at low energies. Integrating out the massive gauge multiplet, these D terms would not appear (corresponding to the fact that effects of small VEV's for $\bar{X}$ would be canceled by shifts of the massive fields). The cancellation of the D terms would arise from the diagram of Fig. 7. In this diagram, the exchange of the massive scalar in the gauge multiplet (one of the superpartners of the massive gauge boson) precisely cancels the quartic couplings associated with the D terms.

When supersymmetry is broken, however, this cancellation is not complete. The scalars in the multiplet are no longer exactly degenerate with the gauge bosons. As a result, there is a quartic coupling remaining in the low-energy theory. Even without a detailed computation, it is easy to determine the sign and order of magnitude of this coupling. Suppose, first that the gauge coupling was zero. Assuming, as we have above, that the F terms all have small VEV's, there will be a nearly degenerate Goldstone supermultiplet, consisting of a Goldstone boson, a light fermion, and an additional scalar particle. This scalar will have a positive mass squared of order $|\lambda\langle F\rangle|^2$, where λ is the coupling to the Goldstino; this is a consequence of a famous sum rule. It is easy to check this in simple examples. In the present case, this leads to a positive quartic coupling of the $\bar{X}$ fields.

Our remarks above can be summarized by the statement that, in addition to the terms in the potential arising from the above superpotential, the potential contains the supersymmetry-breaking terms

$$V_{\text{soft}}=-m_s^2|x^s|^2-m_a^2|x^a|^2+(g_L^2/2)\gamma|x^s|^2|x^a|^2 ,$$

where the last term arises from the incomplete cancellation of the $SU(3)_L$ and $SU(3)_R$ D terms, and is of order $|F|^2/M^4$, where F is a typical F term in the supercolor sector, while M is a typical mass. It is easy to see that, for a range of parameters, the potential has a local minimum at which

$$x^{s2}=3m_s^2/\lambda_4^2,\quad x^a=0,\quad f=\bar{f}=0 . \tag{4.24}$$

What does the theory look like at this minimum? The SU(3) symmetry is still unbroken, but, due to the VEV of x^s, all of the fields which carry SU(3) quantum numbers gain mass. [Note that the gaugino of the SU(3) gains mass also at one loop.[1]] Thus, we have an effective pure SU(3) gauge theory. This R-color theory is now quite asymptotically free, and, depending on the precise values of the $SU(3)_L$ and $SU(3)_R$ couplings, can get strong rather quickly.

[1]Note that, because both the scalar *squared* masses and the gaugino masses arise at one loop, the gauginos are generically lighter than the f, $\bar{f}$, and x^a fields, and will contribute to the renormalization group evolution for a decade or so.

The asymptotic freedom of R color is phenomenologically essential. It gets rid of a Goldstone boson produced by the supercolor interactions. As we have noted earlier, the theory possesses an R symmetry, explicitly broken by $SU(3)_L$ and $SU(3)_R$ anomalies. This symmetry has an ordinary color anomaly and is spontaneously broken by the supercolor sector, giving rise to a not very invisible axion with couplings to ordinary quarks. In order to be consistent with astrophysical bounds, this pseudo Goldstone boson must gain a mass of order a few MeV, at least. The mass of the axion is on the order of Λ_R^2/Λ_7. Since, as we will see, supersymmetry breaking and $SU(2)\times U(1)$ breaking in the ordinary sector are only achieved at two and more loops, we must have the supercolor scale be rather large compared with the weak scale; $\Lambda_7 \sim 10^7$ GeV. Thus the scale Λ_R of R color must be greater than ~ 100 GeV; this in turn requires that the $SU(3)_{L,R}$ couplings be rather large, but there is a finite range of g_L, g_R for which this condition is satisfied, and yet the couplings do not blow up below M_P.

D. The "low-energy" spectrum

We now wish to ask about the spectrum of "ordinary" squarks, sleptons, and gauginos. We will see that, in the effective theory below the scale of the f, $\bar{f}$, and $\bar{X}$ fields, the gauginos gain mass at one loop, while squarks and sleptons gain positive mass-squared at two-loop order. The problem of $SU(2)\times U(1)$ breaking will be taken up in the next section.

We first have to address another potential problem in the model: the appearance of a Fayet-Iliopoulos D term for hypercharge as we integrate out the heavy fields f and $\bar{f}$. Such a term is phenomenologically dangerous since, if it is large, it could lead to very light squarks and sleptons, or even squark VEV's. Suppose $g_L \neq g_R$, and $\lambda_5^t \neq \lambda_5^d$. Then the diagram of Fig. 8 leads, in general, to a nonzero Fayet-Iliopoulos term. The relation $\lambda_5^t = \lambda_5^d$ is renormalized at one loop, and so we assume it does not hold. Thus it is necessary to insist that $g_L = g_R$, to a rather high degree of accuracy (roughly of order α/π). If this were the case, the full theory, ignoring ordinary quarks and leptons, would possess a left-right symmetry which would ensure the absence of a D term. Of course, any such symmetry is broken by the gauge couplings of quarks and leptons, so one must ask how natural the relation $g_L = g_R$ is. First, note that radiative corrections to this relation will arise only at high loop order. Second, recall that in string theory, one typically has equality of various gauge couplings at tree level. If that were the case here, the subsequent evolution of these couplings would induce only small differences in g_L and g_R. Thus it does not seem implausible to make such an assumption.

From now on, we will assume that q_L and q_R are equal, and that any Fayet-Ilipoulos term is very small. We turn to the computation of the gaugino and squark and slepton masses. Again, we consider first the effective theory below the supercolor scale. In this theory, the f and $\bar{f}$ fields have soft supersymmetry-breaking corrections to their masses. If we integrate out these fields, gauginos will obtain mass at one loop from graphs such as those

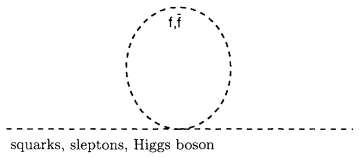

FIG. 8. One-loop diagram contributing to a Fayet-Iliopoulos term in the effective theory.

shown in Fig. 2. These will lead to masses of order

$$m_i = \frac{\alpha_i}{\pi}\frac{\langle F_{\bar{X}}\rangle}{\langle \bar{X}\rangle} . \tag{4.25}$$

Note the result that the gaugino masses are proportional to their gauge couplings squared, just as in the usual grand unified $N=1$ minimal supergravity models. This result only holds when $\langle F_{\bar{X}}\rangle$ is small compared with $\langle \bar{X}\rangle$. Otherwise gaugino masses depend on λ_5^t and λ_5^d; and for $\lambda_5^t \neq \lambda_5^d$ need not satisfy the grand unified theories (GUT) relations.

Squark and slepton masses will arise at two loops from the diagrams shown in Fig. 3. These diagrams are logarithmically divergent. The upper cutoff should be interpreted as the supercolor scale (if one wants to obtain the subleading terms, it is necessary to "open up" the mass insertions, computing three-loop diagrams including supercolor fields). It is not difficult to compute the logarithmic term. [In this computation, it is perhaps worth noting that the separate diagrams exhibit a $\log^2(\Lambda)$ behavior, but the final answer only contains a single logarithm.] One obtains

$$\tilde{m}^2 = -\frac{C_F}{4}\left[\frac{\alpha_i}{\pi}\right]^2 \delta m_f^2 \ln\left[\frac{M^2}{m_f^2}\right] ,$$

where C_F is the quadratic Casimir of the matter representation [e.g., $\frac{4}{3}$ for color triplets, $\frac{3}{4}$ for SU(2) doublets] and δm_f^2 is the supersymmetry-breaking mass shift of the f and $\bar{f}$ fields; note that this quantity is negative.

The main features to note about this result are that it is positive (so color and electric charge can remain unbroken), and that the scalar masses, in this approximation, depend only on gauge quantum numbers, so flavor-changing processes are adequately suppressed. Also, the squark and slepton masses are logarithmically enhanced compared with the gaugino masses.

In summary, the superpartner spectrum in these models is computable, although unfortunately it depends on several new coupling constants. However, assuming that the superpotential couplings $\lambda_5^{t,d}$ are comparable, we can make the following rough predictions for the squark, slepton, and gaugino masses:

$$\begin{aligned} &\frac{m_3}{g_3^2} \approx \frac{m_2}{g_2^2} \approx \frac{3}{5}\frac{m_1}{g_1^2} , \\ &\tilde{m}_q \sim \sqrt{\ln(M^2/m_f^2)}\, m_3 , \\ &\tilde{m}_l \sim (g_2^2/g_3^2)\tilde{m}_q , \\ &\tilde{m}_e \sim (g_1^2/g_3^2)\tilde{m}_q , \end{aligned} \tag{4.26}$$

where $\tilde{m}_q$ is the mass of the (nearly degenerate) squarks, $\tilde{m}_l$ is the mass of the slepton doublets, and $\tilde{m}_e$ is the mass of the slepton singlets. As we will see in the next section, the weak scale is determined by a three-loop negative contribution to the Higgs boson mass squared, which is comparable to $\tilde{m}_l^2$. Thus we expect $\tilde{m}_l \sim v = 250$ GeV, which gives for the approximate size of the other scalar superpartner masses $\tilde{m}_e \sim 100$ GeV, $\tilde{m}_q \sim 900$ GeV. The slepton SU(2) singlets could be within reach of the CERN e^+e^- collider LEP II.

E. SU(2)×U(1) breaking

In this section we turn to the problem of electroweak symmetry breaking. At two loops, we have obtained positive masses for all of the scalar fields in the low-energy theory. If SU(2)×U(1) is to break, at least one Higgs boson field must acquire a negative mass squared. For this to happen, a three-loop negative contribution to the mass squared must be larger than the two-loop contributions. As pointed out long ago [5], in a model such as this one it is easy to obtain a negative mass squared for the Higgs boson which couples to the top quark. The point is that the loop corrections of Fig. 4, while suppressed by a factor $3g_t^2/(16\pi^2)$ are enhanced both by a logarithm of $m_f^2/\tilde{m}_q^2$ and by the fact that the top squark mass itself is proportional to α_s^2 rather than α_W^2, as for the lowest-order Higgs boson mass squared. Thus for top quarks in the presently allowed mass range, this three-loop graph can give a negative contribution to the Higgs boson mass squared which is larger than the two-loop positive contribution. To see the logarithmic enhancement, it is convenient to study the effective theory below the mass scale m_f of the messenger particles. In the low-energy theory, the graph in Fig. 4, which is proportional to the large squark mass squared $\tilde{m}_q^2$ causes the mass squared M_1^2 of H_1 to run. (There will be other contributions to the renormalization-group equations in the effective theory, e.g., from trilinear scalar terms and gaugino masses, but these are smaller.) M_1^2 is positive at m_f, and decreases rapidly. If M_1^2 becomes negative at a scale above $\tilde{m}_q$ then H_1 will get a VEV. However, in order to give masses to all quarks and leptons, both H_1 and H_2 must get VEV's. The symmetries of the model prevent the generation of a $m_{12}^2 H_1 H_2$ term in the potential. Furthermore, it is not easy to obtain a negative mass squared for H_2, since, in general, the bottom quark Yukawa coupling, g_b, is not as large. One can, of course, try to choose couplings so that $\langle H_2 \rangle \ll \langle H_1 \rangle$. In this case the bottom quark Yukawa can be large. However, an examination of the renormalization-group equations shows that this requires a certain amount of fine-tuning (better than 10%). Even if H_2 does obtain a VEV, it is necessary to add additional fields to obtain suitable breaking of SU(2)×U(1). As is well known, in the MSSM, which only has soft supersymmetry breaking, if H_1 and H_2 both obtain negative mass squared, the potential is unbounded below. The present case is somewhat different because not all the supersymmetry-breaking terms induced in the effective theory by radiative corrections are soft; however, the nonsupersymmetric dimension-four terms are much smaller than the supersymmetric terms and do not help give an acceptable symmetry breaking. Moreover, in the absence of an $H_1 H_2$ coupling, the theory has a Peccei-Quinn symmetry and one obtains a standard axion. To get around this we add a singlet field, with couplings

$$W_S = \tilde{\lambda}_1 S H_1 H_2 + (\tilde{\lambda}_2/3) S^3 . \tag{4.27}$$

In order to understand the absence of other terms, one can invoke a discrete symmetry. The terms in (4.27) gives rise to an effective quartic coupling of the Higgs boson fields, which prevents the runaway behavior. So one might hope that with this modification, with a negative mass squared only for H_1, we could obtain a sensible breaking of SU(2)×U(1).

The S scalar will obtain a mass at one higher order in the loop expansion than the Higgs boson fields, and trilinear terms involving the scalar are also of higher order. So, to get a feeling for what may happen, we simply examine the potential

$$\begin{aligned} V = &-m^2|H_1|^2 + m'^2|H_2|^2 + (g^2/8)(H_1^\dagger \tau^a H_1 + H_2^\dagger \tau^a H_2)^2 \\ &+ (g'^2/8)(H_1^\dagger H_1 - H_2^\dagger H_2)^2 + |\tilde{\lambda}_1|^2(|H_1 S|^2 + |H_2 S A|^2) \\ &+ |\tilde{\lambda}_1 \epsilon_{ij} H_{1i} H_{2j} + \tilde{\lambda}_2 S^2|^2 . \end{aligned} \tag{4.28}$$

However, a detailed study of this potential shows that there is always a scalar field with a mass less than about 40 GeV. In fact, as it stands, this potential possesses a global symmetry which leads to a massless pseudoscalar. Corrections to the potential, such as the nonsupersymmetric cubic terms, will break this symmetry, and give the light pseudoscalar a mass of order a few GeV.

This situation is unacceptable, given the strong LEP limits on the decay $Z \to$ scalar + pseudoscalar. Preliminary estimates of further radiative effects indicate that these will not help much. So we need to consider some further modification. The simplest possibility seems to be to add a set of vectorlike quarks. In order to maintain the successful grand unification of the SU(3)×SU(2)×U(1) coupling constants, we can also add vectorlike leptons. We take these fields to have the quantum numbers of a 5 and $\bar{5}$ of ordinary SU(5). Denote the corresponding quarks by D and $\bar{D}$. These can couple to the S field, $SD\bar{D}$, and the $\bar{D}$ field can couple to the Higgs field, H_2, $H_2 q \bar{D}$, where q is an ordinary quark doublet field. If these additional couplings are large enough, several things will happen. First, the S field will also obtain a large negative mass squared at one loop. It can thus obtain a large expectation value, giving rise to a mass for the fermionic components of the D and $\bar{D}$ fields (and the corresponding leptons). Second, the field H_2 could obtain a large negative contribution to its mass squared at one loop. Also, the large S VEV, in conjunction with the VEV of H_1, will induce a VEV for H_2. So with this modification, a sensible breaking of SU(2)×U(1) can arise. The parameter space of this model is quite large, and we will not attempt a complete exploration here. However, it is clear that there are finite ranges of parameters for which a sensible spectrum is obtained, with all the scalars heavier than the Z. There is generi-

cally a light pseudoscalar, with mass in the GeV range, which is mainly a gauge singlet. Note that with this minimal set of extra fields, QCD is no longer asymptotically free above the scale of the f and $\bar{f}$ fields. However, it is almost asymptotically free, and does not hit its Landau pole until extremely high energy.

One potential problem with this scenario is that the charge $\frac{1}{3}$ quarks in q will mix with those in D, giving rise to flavor-changing neutral currents (FCNC's) involving the Z. However, S receives no positive two-loop contribution to its mass squared, and will get a larger VEV than $H_{1,2}$. Thus the D will be heavier than the weak scale, which greatly suppresses its mixing with the light d and s quarks. The necessary suppression of the most severe FCNC, $K_L \to \mu^+\mu^-$, is easily achieved by requiring that all superpotential couplings of the lightest two families be smaller than about 10^{-1}, including couplings of the form $H_2 q_{1,2}\bar{D}$. This assumption is consistent with the small masses for these two families and is natural, in the sense of 't Hooft [11], since a chiral flavor symmetry is restored in the limit that all couplings of the lightest quark flavors vanish.

F. Conclusions

We have seen that it is, in fact, possible to construct models with dynamical supersymmetry breaking at relatively low energy. We have exhibited a model in which the following occurs.

(a) $SU(2)\times U(1)$ is properly broken.

(b) All superpartners have adequate, calculable masses.

(c) There is enough degeneracy among quark and lepton masses to assure absence of flavor-changing neutral currents. This occurs naturally as a result of the accidental flavor symmetry of the gauge interactions.

(d) There is no new source of CP violation in the low-energy theory, explaining the absence of large particle electric dipole moments.

(e) There are no dangerous axions or Goldstone bosons.

(f) All couplings are small up to very high energies.

(g) It is still possible to unify $SU(3)\times SU(2)\times U(1)$.

(h) The superpotential is the most general cubic potential allowed by the gauge symmetries and is the most general consistent with a set of (anomalous) discrete symmetries.

(i) The gravitino is light, of order a keV, and hence provides no cosmological problems [12].

The model we have described should be viewed as an existence proof. Probably the most serious drawback of this particular model is that the potentially dangerous axion only gains adequate mass if a certain gauge coupling is in a particular range (the lower limit set by the mass; the upper limit set by the requirement that the gauge coupling not blow up too soon). It would be nice to find a more natural model which does not suffer from this difficulty.

Some features of the model appear to be generic. First, the squark and slepton masses are, to a good approximation functions only of their gauge quantum numbers. Second, the need for additional fields in order to break $SU(2)\times U(1)$ is almost certainly general. The choice we have described here, of an additional singlet as well as a set of vectorlike fermions, is the simplest possibility we have found.

Finally, we would like to comment on some cosmological issues and problems with this model. Perhaps the most serious potential problem is one of domain walls. The model possesses several spontaneously broken discrete symmetries. Fortunately, all of them possess anomalies with respect to one of the strong gauge groups. The corresponding domain walls will thus disappear by the mechanism of Preskill *et al.* [13].

In general, however, one might expect nonanomalous discrete symmetries to arise (this occurs, for example, in the model without the mirrors). However, it is not clear that the problem is severe. Indeed, the clue to a solution lies in the solution of Preskill *et al.* These authors noted that, even if the scale of the spontaneous breaking is 10's of GeV, the tiny lifting of the degeneracy (12 orders of magnitude smaller) by QCD is enough to cause collapse of the walls, simply because the expansion of the Universe is so slow. Suppose, then, that one has some discrete symmetry without anomalies in the low-energy theory. If this summary is broken by nonrenormalizable terms, this will lead to a breaking of the degeneracy. Even if this effect is quite small, it can be sufficient to get rid of the domain walls. For example, dimension-five operators with coefficients slightly larger than $1/M_P$ or dimension-six operators associated with a scale of order 10^9 GeV or smaller should be enough. This solution to the domain-wall problem can be relevant quite generally, and has antecedents [14] in earlier work on axions and technicolor.[2]

In these models, one must also study the possibility of stable or nearly stable massive particles, such as the f fermions. Again, it may be necessary to invoke higher dimension operators to allow these to decay and avoid cosmological problems. This problem may not be generic, but specific to the model under study.

In any case, we believe that models with low-energy dynamical supersymmetry breaking in the visible sector are a plausible alternative to more conventional hidden sector supergravity models. They solve some of the most troubling problems of the hidden sector models, and they provide a dynamical solution of the hierarchy problem.

ACKNOWLEDGMENTS

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[2]Recently it has been discussed for spontaneous breaking of P and CP symmetries [15]. However, in the cases which have been studied, there are low dimension operators in the low-energy theory which can break the symmetry. Moreover, the must plausible context for spontaneous breaking of these symmetries is in theories in which these symmetries are gauge symmetries, in which case there is no explicit breaking of the symmetry [16].

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