14 Dynamical supersymmetry breaking

One of the original reasons for the interest in supersymmetry was the possibility of dynamical supersymmetry breaking. So far, however, we have exhibited models in which supersymmetry is unbroken in the true ground state, as in the case of QCD with only massive quarks or models with moduli spaces or approximate moduli spaces. In this chapter, we describe a number of models in which a non-trivial dynamics breaks supersymmetry. We will see that dynamical supersymmetry breaking occurs under special, but readily understood, conditions. In some cases we will be able to exhibit this breaking explicitly, through systematic calculations. In others we will have to invoke more general arguments. Then we will turn to theories in which supersymmetry is preserved in the lowest energy state but in which there exist metastable states with broken supersymmetry. We will argue that this is a generic phenomenon and see that it is even sometimes true in massive QCD.

14.1 Models of dynamical supersymmetry breaking

We might ask why, so far, we have not found supersymmetry to be dynamically broken. In supersymmetric QCD with massive quarks, we might give the Witten index as an explanation. We might also note that there is no promising candidate for a goldstino. With massless quarks we have flat directions and, as the fields get larger, the theory becomes more weakly coupled so that any potential tends to zero.

This suggests two criteria for finding models with dynamical supersymmetry breaking (DSB).

1. The theory should have no flat directions at the classical level.
2. The theory should have a spontaneously broken global symmetry.

The second criterion implies the existence of a Goldstone boson. If the supersymmetry were unbroken, any would-be Goldstone boson must lie in a multiplet with another scalar as well as a Weyl fermion. This other scalar, like the Goldstone particle, has no potential so the theory has a flat direction. But, by assumption, the theory classically (and therefore almost certainly quantum mechanically) has no flat direction. So supersymmetry is likely to be broken. These criteria are heuristic but, in practice, when a systematic analysis is possible, they always turn out to be correct.

Perhaps the simplest model with these features is a supersymmetric $SU(5)$ theory with a single $\bar{5}$ and a single 10 representation. In the exercises, you can show that this theory, in fact, has no flat directions and that it has two non-anomalous $U(1)$ symmetries. One can

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give arguments showing that these symmetries are broken. So it is likely that this theory breaks supersymmetry.

However, this is a strongly coupled model and it is difficult actually to prove that supersymmetry is broken. In the next section, we will describe a simple weakly coupled theory in which dynamical supersymmetry breaking occurs within a controlled approximation.

14.1.1 The $(3, 2)$ model

A model in which supersymmetry turns out to be broken is the (3, 2) model. This theory has gauge symmetry $SU(3) \times SU(2)$, and matter content

$$Q(3,2), \quad \bar{U}(\bar{3},1), \quad L(1,2), \quad \bar{D}(\bar{3},1). \tag{14.1}$$

This is similar to the field content of a single generation of the standard model, but without the extra $U(1)$ and the positron. The most general renormalizable superpotential consistent with the symmetries is

$$W = \lambda Q L \bar{U}. \tag{14.2}$$

This model admits an R symmetry that is free of anomalies. There is also a conventional $U(1)$ symmetry, under which the charges of the various fields are the same as in the standard model (one can gauge this symmetry if one also adds an e^+ field).

While this model has global symmetries, it is different from supersymmetric QCD in that it does not have classical flat directions. To see this, note that by $SU(3) \times SU(2)$ transformations one can bring Q to the form

$$Q = \begin{pmatrix} a & 0 \\ 0 & b \\ 0 & 0 \end{pmatrix}. \tag{14.3}$$

Now, the vanishing of the $SU(2)$ D term forces

$$L = \left(0, \sqrt{|a^2| - |b^2|}\right). \tag{14.4}$$

The vanishing of the F terms for $\bar{u}$ requires $|a| = |b|$. Then the vanishing of the $SU(3)$ D term forces

$$\bar{U} = \begin{pmatrix} a' \\ 0 \\ 0 \end{pmatrix}, \quad \bar{D} = \begin{pmatrix} 0 \\ a'' \\ 0 \end{pmatrix} \tag{14.5}$$

(up to interchange of the two vevs), with

$$|a'| = |a''| = |a|.$$

Finally, the $\partial W/\partial L$ equations lead to $a = 0$.

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To analyze the dynamics of this theory, consider first the case where $\Lambda_3 \gg \Lambda_2$. Ignoring, at first, the superpotential term this is just $SU(3)$ with two flavors. In the flat direction of the D terms there is a non-perturbative superpotential

$$W_{\rm np} = \frac{\Lambda^5}{\det \overline{Q}Q} \sim \frac{1}{v^4}. \tag{14.6}$$

The full superpotential in the low-energy theory is the sum of this term and the perturbative term. It is straightforward to minimize the potential and establish that supersymmetry is broken. One finds

$$a = 1.287\frac{\Lambda}{\lambda^{1/7}}, \quad b = 1.249\frac{\Lambda}{\lambda^{1/7}}, \quad E = 3.593\lambda^{10/7}\Lambda^4. \tag{14.7}$$

If $\Lambda_2 \gg \Lambda_3$, supersymmetry is still broken but the mechanism is different. In this case, before including the classical superpotential the strongly coupled theory is $SU(2)$ with two flavors. This is an example of a model with a *quantum moduli space*. This notion will be explained in the next chapter but it implies that $\langle QL \rangle \neq 0$, so at low energies there is a superpotential (F term) for $\bar{U}$.

There does not exist, at the present time, an algorithm to generate all models which exhibit dynamical supersymmetry breaking, but many classes have been identified. A generalization of the $SU(5)$ model, for example, is provided by an $SU(N)$ model with an antisymmetric tensor field A_{ij} and $N-4$ $\bar{F}$ terms. It is also necessary to include a superpotential,

$$W = \lambda_{ab} A \bar{F}^a \bar{F}^b. \tag{14.8}$$

Other broad classes are known, including generalizations of the $(3, 2)$ model. A somewhat different, and particularly interesting, set of models is described in Section 15.4. Catalogs of known models, as well as studies of their dynamics, are given in some references in the suggested reading at the end of this chapter.

14.2 Metastable supersymmetry breaking

In the previous section we established criteria for dynamical supersymmetry breaking and exhibited an example, the $(3, 2)$ model, which satisfies the criteria and exhibits dynamic supersymmetry in a stable ground state. But there are a number of ways in which we might view these criteria as limiting. First, while there are many models which satisfy them, they seem exceptional and not particularly generic. Second, it is difficult to build realistic models without spoiling the chiral structure of these theories. Finally, the criteria themselves are troubling, especially the requirement of a continuous global symmetry. We do not expect such symmetries in theories of gravity, so these symmetries must arise as accidents and must hold to some high degree of accuracy. Indeed, these criteria seem less sharp in the framework of supergravity.

If we consider theories with *metastable* ground states, i.e. theories having a stable ground state with unbroken supersymmetry but where supersymmetry is broken in a higher-energy, classically stable, state, the possibilities are greatly enlarged. Indeed, we can consider this

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question in the O'Raifeartaigh models. Rather than imposing a continuous R symmetry, we can consider discrete symmetries, for example a Z_N subgroup of a continuous R symmetry. For the fields $Z,\ Y,\ A$ we can require, with $\alpha = e^{\frac{2\pi i}{N}}$,

$$Z \to \alpha^2 Z, \quad Y \to \alpha^2 Y, \quad A \to A \tag{14.9}$$

while the superpotential transforms as

$$W \to \alpha^2 W. \tag{14.10}$$

Imposing, for simplicity, an additional symmetry $A \to -A,\ Y \to -Y$, the most general *renormalizable* superpotential takes the form of a simple O'Raifeartaigh model but where, beyond the renormalizable level, additional couplings are allowed:

$$W = Z(A^2 - \mu^2) + mYA + \frac{Z^{N+2}}{M^{N-1}} + \cdots . \tag{14.11}$$

Focusing just on the Z^{N+2} term, there is now a supersymmetric vacuum at

$$(N+2)Z^{N+1} = \mu^2 M^{N-1}. \tag{14.12}$$

For M large (e.g of order the Planck or unification scale) compared with μ this vacuum is far away. Near the origin, the Coleman–Weinberg calculation still leads to a local minimum of the potential. The time required to tunnel from the metastable vacuum to the supersymmetric vacuum grows exponentially with power M/μ (on including effects of general relativity, the time often becomes infinite). So this instability is not a phenomenological concern.

One might imagine that the phenomenon of metastable supersymmetry breaking in theories with discrete R symmetries is rather generic. In models with singlet chiral fields and a continuous R symmetry, if all fields have R charge 0 or 2 then supersymmetry breaking occurs when the number of fields X_i with charge 2 exceeds the number A_α with charge 0. A similar statement holds for the discrete symmetries.

14.2.1 Metastable dynamical supersymmetry breaking: the ISS model

The phenomenon of dynamical metastable supersymmetry breaking appears, then, to be rather generic. Remarkably, this already occurs in supersymmetric QCD with $N_\mathrm{f} > N_\mathrm{c}$, with *massive* quarks, as first pointed out by Intriligator, Shih and Seiberg (ISS). We have already explained that, quite generally, supersymmetric theories with massive, vector-like, fields do not break supersymmetry, in the sense that they possess multiple (typically N, for the gauge group $SU(N)$) supersymmetric ground states. But, consider the case $3N/2 > N_\mathrm{f} \geq N_\mathrm{c} + 1$. Turning off the mass term we will see in Section 16.4 that the theory is dual to a theory with gauge group $SU(N_\mathrm{f} - N_\mathrm{c})$, with:

1. N_f quarks in the fundamental representation, q_f, transforming in the $(1, N_\mathrm{f})$ representation of the flavor symmetry, $SU(N_\mathrm{f})_\mathrm{L} \times SU(N_\mathrm{f})_\mathrm{R}$;
2. N_f in the antifundamental representation, transforming as $(\bar{N}_\mathrm{f}, 1)$ under flavor;
3. a chiral field $\Phi_{f,\bar{f}}$, transforming in the $(N_\mathrm{f}, \bar{N}_\mathrm{f})$ representation, which is a singlet of the dual gauge group.

The superpotential of the magnetic theory is

$$W_{\text{mag}} = \mu \bar{q} \Phi q. \tag{14.13}$$

Now turn on a small mass term in the underlying, "electric", theory,

$$\delta W = \bar{Q} m Q. \tag{14.14}$$

We expect the appearance of a small term proportional to m, in the dual, "magnetic", theory. The term $\mu \,\mathrm{Tr}\, \Phi \; m$ transforms under the global symmetries (including the anomalous $U(1)$s) in the same way as the original mass term. So we will assume that it is in fact present, i.e. that the full superpotential of the magnetic theory is

$$W_{\text{mag}} = \mu \bar{q} \Phi q + \mu \,\mathrm{Tr}\, \Phi \, m. \tag{14.15}$$

Recalling that the fields q, $\bar{q}$ are fundamentals of the dual gauge group and requiring that the D term conditions of this group be satisfied, the vacuum of the dual theory breaks supersymmetry. It is important that $N_{\text{f}} - N < N_{\text{f}}$; the resulting breaking is called "rank breaking". One can see this by using the flavor symmetries to write, for example, for $N_{\text{c}} = 2$, and $N_{\text{f}} = 3$,

$$q = \bar{q} = \begin{pmatrix} v_1 & 0 & 0 & \cdots \\ 0 & v_2 & 0 & \cdots \\ 0 & 0 & v_3 & \cdots \end{pmatrix}. \tag{14.16}$$

With this choice, we can satisfy the equations

$$\frac{\partial W}{\partial \Phi_{f,\bar{f}}} = 0 \tag{14.17}$$

only for $f, \bar{f} = 1, 2, 3$, not for larger f. This generalizes to the other values of $N_{\text{f}}, N_{\text{c}}$ in this class of models.

It still remains to verify that there is a good non-supersymmetric vacuum in the magnetic theory. For this, we need to consider the pseudomoduli of the classical theory. These are components of Φ, essentially those components which cannot gain mass by mixing with the $q_f, \bar{q}_f$ superfields. Clearly, in particular the components of $\Phi_{f,\bar{f}}$ with $f, \bar{f} > N$ are massless at the tree level. A Coleman–Weinberg calculation is necessary to determine the masses of these fields and to establish whether $\Phi = 0$ is a good ground state. The answer turns out to be yes.

We know that in the electric theory there are N supersymmetric ground states. These can be found in the magnetic description; decays to them are highly suppressed for small quark mass.

14.2.2 Retrofitting

A broad class of models exhibiting dynamical metastable supersymmetry breaking can be found by starting with the O'Raifeartaigh models. Again, a simple example is that of Eq. (14.11) above. Now, however, we replace the dimensional parameters m and μ^2 by couplings to a strongly interacting group which generates these scales dynamically.

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For simplicity, we will consider μ^2. We introduce an $SU(N)$ gauge group with field strength W_α,

$$W = \lambda Z A^2 + \frac{Z}{M} W_\alpha^2 + mYA \tag{14.18}$$

(we will see that couplings of chiral fields to gauge fields of this type are common in string theory, where M might be the Planck scale or the scale of the string theory). Gaugino condensation in the $SU(N)$ group gives rise to an expectation value Λ^3 for W_α^2,

$$W = Z(A^2 - \mu^2 e^{-Z/(Mb_0)}) + mYA, \tag{14.19}$$

where b_0 is the beta function of the gauge theory.

Near the origin the Coleman–Weinberg calculation is identical to that of the O'Raifeartaigh model, and the potential has a minimum at $Z = 0$. But clearly there are lower energy states at larger fields due to:

1. the exponential term in Eq. (14.19);
2. possible higher-order terms in powers of Z/M_p.

Models of this type illustrate the fact that metastable dynamical supersymmetry breaking is a generic phenomenon in supersymmetric field theories. They vastly expand the possibilities for supersymmetric model building.

We have seen, in this section, that the dynamical breaking of supersymmetry is common. Flat directions are often lifted and, in many instances the supersymmetry is broken with a stable ground state. So, we are ready to address the question: how might supersymmetry be broken in the real world?

14.3 Particle physics and dynamical supersymmetry breaking

14.3.1 Gravity mediation and dynamical supersymmetry breaking: anomaly mediation

One simple approach to model building which we explored in Chapter 11 was to treat a theory which breaks supersymmetry as a "hidden sector". This construction, as we presented it, was rather artificial. If we replace, say, the Polonyi sector by a sector which breaks supersymmetry dynamically, the situation is dramatically improved. If we suppose that there are some fields transforming under only the Standard Model gauge group and some transforming under only the gauge group responsible for symmetry breaking, the visible/hidden sector division is automatic. As we will see, this sort of division can arise rather naturally in string theory.

In such an approach the scale of supersymmetry breaking is again $m_{3/2}M_\mathrm{p}$, where we now understand this scale as the exponential of a small coupling at a high-energy scale (presumably the Planck, GUT or string scale). For scalars, soft-supersymmetry-breaking

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masses and couplings arise just as they did previously. There is no symmetry reason why these masses should exhibit any sort of universality.

One puzzle in this scenario is related to gluino masses. Examining the supergravity Lagrangian, the only term which can lead to gaugino masses is

$$\mathcal{L}_{\lambda\lambda} = e^{K/2} f'_{\alpha\beta k}(D_k W)\lambda^\alpha \lambda^\beta. \tag{14.20}$$

Here f is the gauge coupling function. So, in order to obtain a substantial gaugino mass, it is necessary that there be gauge-singlet fields with non-zero F terms. In most models of stable dynamical supersymmetry breaking there are no scalars which are singlets under all the gauge interactions. In metastable models, such as retrofitted models, it is necessary to suppose that there is some sort of discrete symmetry which accounts for the absence of certain couplings. These symmetries will forbid the coupling of hidden sector fields to visible sector gauge fields through low-dimension operators. In other words, we do not have couplings of the form

$$\frac{S}{M} W_\alpha^2, \tag{14.21}$$

where the F component of S has a non-zero vev. This suggests that gaugino masses would be suppressed relative to squark and slepton masses by powers of $M_{\rm int}/M_{\rm p}$.

But this turns out to be not quite correct. This is associated with a phenomenon known as "anomaly mediation". The term is arguably a misnomer; no actual symmetry of the theory is anomalous. The appearance of these terms can be understood, in some cases, as an issue of locality: the gaugino masses are themselves local but the supersymmetric operator which gives rise to them is not (i.e. it includes non-local terms). In other cases, a completely Wilsonian description is not available. Here we simply note that such terms are, in many instances, *required* by supersymmetry. Consider for example a pure $SU(N)$ gauge theory coupled to supergravity, with a small constant W_0 in the superpotential. In this theory, gaugino condensation occurs and gives rise to a non-perturbative correction to the superpotential,

$$W_{\rm np} = -\frac{N}{32\pi^2}\langle\lambda\lambda\rangle.$$

From $V = -3|W_0 + W_{\rm np}|^2$, then, we predict the following term in the potential:

$$-\frac{3N}{32\pi^2} W_0^* \langle\lambda\lambda\rangle. \tag{14.22}$$

It is natural to interpret this as resulting from an underlying term in the action,

$$\delta\mathcal{L} = -\frac{b_0}{16\pi^2} m_{3/2}\lambda\lambda. \tag{14.23}$$

One can argue for the presence of such a term for all N and $N_{\rm f}$ in a similar fashion. But the term can be found more directly from the structure of the underlying supergravity theory.

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14.3.1.1 Split supersymmetry

The anomaly-mediated expression for the gaugino masses suggests an approach to model building of particular interest, given the large mass scale for squarks suggested by the Higgs mass. Even if one is willing to accept some fine tuning, one might need lighter gauginos to account for WIMP dark matter and to improve the quality of gauge coupling unification. If X denotes the field, with a non-vanishing F component, responsible for supersymmetry breaking, then one might suppose that there is no XW_α^2 coupling. In this case, assuming that the scalar masses are of order $m_{3/2}$, one can contemplate gauginos with masses lighter by a loop factor. So, for example, if squarks are at 30 TeV, one might have gluinos at scales slightly above one TeV and winos (the LSP), according to Eq. (14.23), a factor 3 or so lighter. One can debate how generic a phenomenon this might be.

14.3.2 Low-energy dynamical supersymmetry breaking: gauge mediation

An alternative to the conventional supergravity approach is to suppose that supersymmetry is broken at some much lower energy, with gauge interactions serving as the messengers of supersymmetry breaking. The basic idea is simple. One again supposes that one has some set of new fields and interactions which break supersymmetry. Some of these fields are taken to carry ordinary Standard Model quantum numbers, so that "ordinary" squarks, sleptons and gauginos can couple to them through gauge loops. This approach, which is referred to as *gauge mediated supersymmetry breaking* (GMSB), has a number of virtues.

1. It is highly predictive: as few as two parameters describe all soft breakings.
2. The degeneracies required to suppress flavor-changing neutral currents are automatic.
3. GMSB easily incorporates DSB and so can readily explain the hierarchy.
4. GMSB makes dramatic and distinctive experimental predictions.

The approach, however, also has drawbacks. Perhaps most serious is related to the "μ problem", which we discussed in the context of the MSSM. In theories with high-scale supersymmetry breaking we saw that there is not really a problem at all; a μ term of order the weak scale is quite natural. The μ problem, however, finds a home in the framework of low-energy breaking. The difficulty is that if one is trying to explain the weak scale dynamically then one does not want to introduce the μ term by hand. Various solutions have been offered for this problem. One possibility is that it is protected by symmetries and generated by the same dynamics which generates supersymmetry breaking. In the rest of our discussion we will simply assume that a μ term has been generated in the effective theory and will not worry about its origin.

14.3.2.1 Minimal gauge mediation (MGM)

The simplest model of gauge mediation contains, as messengers, a vector-like set of quarks and leptons, q, $\bar{q}$, ℓ and $\bar{\ell}$. These have the quantum numbers of a 5 and a $\bar{5}$ representation of $SU(5)$. The superpotential is taken to be

$$W_{\rm mgm} = \lambda_1 q\bar{q} + \lambda_2 S\ell\bar{\ell}. \qquad (14.24)$$

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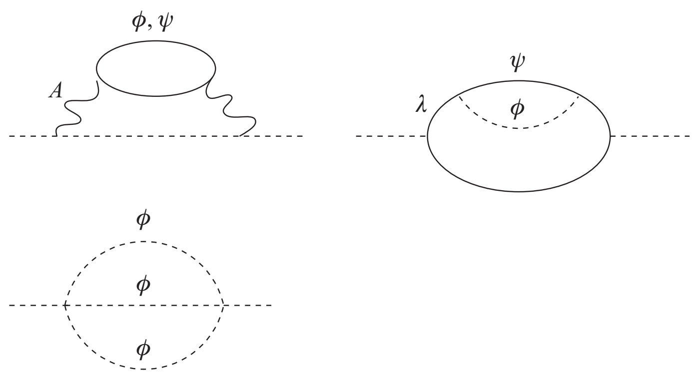

Fig. 14.1 Two-loop diagrams contributing to squark masses in a simple model of gauge mediation.

We will suppose that some dynamics gives rise to non-zero expectation values for S and F_S. We will not provide here a complete microscopic model to explain the origin of the parameters F_S and $\langle S \rangle$ that will figure in our subsequent analysis; retrofitting provides one strategy. To find a compelling model of the underlying dynamics is a good research problem. Instead, we will go ahead and immediately compute the superparticle spectrum for such a model. Ordinary squarks and sleptons gain mass through the two-loop diagrams shown in Fig. 14.1. While the prospect of computing a set of two-loop diagrams may seem intimidating, the computation is actually quite easy. If one treats F_S/S as small then there is only one scale in the integrals. It is a straightforward matter to write down the diagrams, introduce Feynman parameters and perform the calculation. There are also various non-trivial checks. For example, the sum of the diagrams must vanish in the supersymmetric limit. These masses can alternatively be computed by writing down an effective action in terms of spurion fields and computing the wave function renormalization factors as functions of the spurions.

One obtains the following expressions for the scalar masses:

$$\widetilde{m}^2 = 2\Lambda^2 \left[C_3 \left(\frac{\alpha_3}{4\pi} \right)^2 + C_2 \left(\frac{\alpha_2}{4\pi} \right)^2 + \frac{5}{3} \left(\frac{Y}{2} \right)^2 \left(\frac{\alpha_1}{4\pi} \right)^2 \right], \tag{14.25}$$

where $\Lambda = F_S/S$, $C_3 = 4/3$ for color triplets and zero for singlets and $C_2 = 3/4$ for weak doublets and zero for singlets. For the gaugino masses one obtains

$$m_{\lambda_i} = C_i \frac{\alpha_i}{4\pi} \Lambda. \tag{14.26}$$

This expression is valid only to lowest order in Λ. Higher-order corrections have been computed; it is straightforward to compute them exactly in Λ.

All these masses are positive and they are described in terms of a single new parameter, Λ. The lightest new particles are the partners of the $SU(3) \times SU(2)$ singlet leptons. If their masses are of order 100 GeV, we have that $\Lambda \sim 30$ TeV. The spectrum has a high degree of degeneracy. In this approximation the masses of the squarks and sleptons are functions only of their gauge quantum numbers, so flavor-changing processes are suppressed.

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Flavor violation arises only through Yukawa couplings, and these can appear only in graphs at high loop order; it is further suppressed because all but the top Yukawa coupling is small.

Apart from the parameter Λ, one has the μ and B_μ parameters (B_μ is the coefficient of the soft-breaking $H_U H_D$ term in the potential; μ and B_μ are both complex), for a total of five. This is three beyond the minimal Standard Model. If the underlying susy-breaking theory conserves CP, this can eliminate the phases, reducing the number of parameters by two.

14.3.2.2 $SU(2)\times U(1)$ breaking

At lowest order, all the squark and slepton masses are positive. The large top quark Yukawa coupling leads to large corrections to $m^2_{H_U}$, however, which tend to drive it negative. The calculation is just a repeat of the one we did in the case of the MSSM. Treating the mass of $\tilde{t}$ as independent of momentum is consistent provided that we cut the integral off at a scale of order Λ (at this scale the calculation leading to Eq. (14.25) breaks down, and the propagator falls rapidly with momentum) and we have

$$m^2_{H_U} = \left(m^2_{H_U}\right)_0 - \frac{6y_t^2}{16\pi^2}\ln\frac{\Lambda^2}{\widetilde{m}_t^2}\left(\widetilde{m}_t^2\right)_0. \tag{14.27}$$

While the loop correction is nominally three-loop in order, because the stop mass arises from gluon loops while the Higgs mass arises at lowest order from W loops we have a substantial effect,

$$\left(\frac{\widetilde{m}_t^2}{m^2_{H_U}}\right)_0 = \frac{16}{9}\left(\frac{\alpha_3}{\alpha_2}\right)^2 \sim 20 \tag{14.28}$$

and the Higgs mass-squared is negative. These contributions are quite large and, given the large value of the Higgs mass, it is again necessary to tune the μ term and other possible contributions to the Higgs mass to a high degree in order to obtain sufficiently small W and Z masses.

14.3.2.3 General gauge mediation

The minimal model of gauge mediation of the previous section makes a quite sharp set of predictions. These predictions, in fact, are referred to as *minimal gauge mediation* (MGM). It is clearly of interest to ask how general they are. It turns out that they are peculiar to our assumption that there is a single set of messengers and that just one singlet is responsible for supersymmetry breaking and R symmetry breaking. Indeed, our messengers have the quantum numbers of a 5 and a $\bar{5}$ representation of $SU(5)$. If, for example, we had considered two singlets, Z_1 and Z_2, with Z_i and F_i non-zero, we could have obtained independent soft-breaking masses for squarks and leptons. Had we allowed different singlets, and taken a 10 and $\overline{10}$ for the messengers, we could have obtained a richer spectrum. Meade, Seiberg and Shih formulated the problem of gauge mediation in a general way and dubbed this formulation *general gauge mediation* (GGM).

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They studied the problem in terms of the correlation functions of (gauge) supercurrents. Analyzing the restrictions imposed by Lorentz invariance and supersymmetry on these correlation functions, they found that the general gauge-mediated spectrum is described by three complex parameters and three real parameters. The spectrum can be significantly different from that of the MGM, but the masses are still only functions of the gauge quantum numbers and flavor problems are still mitigated.

The basic structure of the spectrum is readily described. In the formulas for the fermion masses we introduce a separate complex parameter m_i, $i = 1, 2, 3$ for each Majorana gaugino. Similarly, for the scalars we introduce a real parameter Λ_c^2 for the contributions from $SU(3)$ gauge fields Λ_w^2 for those from $SU(2)$ gauge fields and Λ_Y^2 for those from hypercharge gauge fields:

$$\widetilde{m}^2 = 2\left[C_3\left(\frac{\alpha_3}{4\pi}\right)^2\Lambda_c^2 + C_2\left(\frac{\alpha_2}{4\pi}\right)^2\Lambda_w^2 + \frac{5}{3}\left(\frac{Y}{2}\right)^2\left(\frac{\alpha_1}{4\pi}\right)^2\Lambda_Y^2\right]. \tag{14.29}$$

One can construct models which exhibit the full set of parameters. In MGM the messengers of each set of quantum numbers each have a supersymmetric contribution to their masses, λM, while the supersymmetry-breaking contribution to the scalar masses goes as λM^2, so in the ratio of these two contributions the coupling cancels out. In GGM model building, additional fields and couplings lead to more complicated relations.

One feature of MGM which is not immediately inherited by GGM is the suppression of new sources of CP violation. Because the gaugino masses are independent parameters, in particular, they introduce additional phases which are inherently CP-violating. Providing a natural explanation of the suppression of these phases is one of the main challenges of GGM model building.

14.3.2.4 Light gravitino phenomenology

There are other striking features of these models. One of the more interesting is that the lightest supersymmetric particle, or LSP, is the gravitino. Its mass is

$$m_{3/2} = 2.5\left(\frac{F}{(100\ \text{TeV})^2}\right)\text{eV}. \tag{14.30}$$

The next-to-lightest supersymmetric particle, or NLSP, can be a neutralino or a charged right-handed slepton. The NLSP will decay to its superpartner plus a gravitino in a time long compared with typical microscopic times but still quite short. The lifetime can be determined from low-energy theorems, in a manner reminiscent of the calculation of the pion lifetime. Just as the chiral currents are linear in the (nearly massless) pion field,

$$j^{\mu 5} = f_\pi \partial^\mu \pi, \quad \partial_\mu j^{\mu 5} = \partial^2 \pi \approx 0, \tag{14.31}$$

so the supersymmetry current is linear in the goldstino G,

$$j_\alpha^\mu = F\gamma^\mu G + \sigma^{\mu\nu}\lambda F_{\mu\nu} + \cdots, \tag{14.32}$$

where F, here, is the goldstino decay constant. From this, if one assumes that the LSP is mostly photino then one can calculate the amplitude for $\tilde{\gamma} \to G + \gamma$ in much the same

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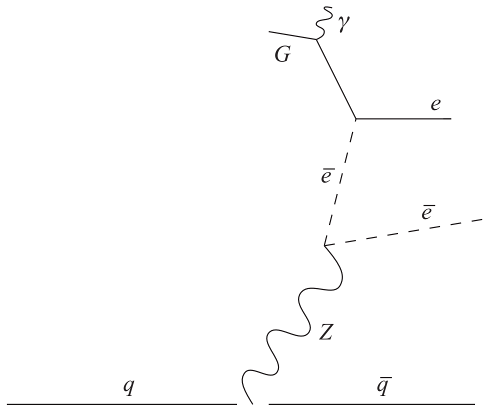

Fig. 14.2 Decay leading to $e^+e^-\gamma\gamma$ events.

way as one considers processes in current algebra. From Eq. (14.32) one sees that $\partial_\mu j^\mu_\alpha$ is an interpolating field for G, so

$$\langle G\gamma|\tilde{\gamma}\rangle = \frac{1}{F}\langle\gamma|\partial_\mu j^\mu_\alpha|\tilde{\gamma}\rangle. \tag{14.33}$$

The matrix element can be evaluated by examining the second term in the current, Eq. (14.32), and noting that $\partial\!\!\!/ \chi = m_\lambda\lambda$.

Given the matrix element, the calculation of the NLSP lifetime is straightforward and yields

$$\Gamma(\tilde{\gamma} \to G\gamma) = \frac{\cos^2\theta_W\, m^5_{\tilde{\gamma}}}{16\pi F^2}. \tag{14.34}$$

This yields a *decay length*:

$$c\tau = 130\left(\frac{100\ \text{GeV}}{m_{\tilde{\gamma}}}\right)^5\left(\frac{\sqrt{F}}{100\ \text{TeV}}\right)^4 \mu\text{m}. \tag{14.35}$$

In other words, if F is not too large then the NLSP may decay in the detector. One even has the possibility of measurable displaced vertices. The signatures of such low decay constants would be quite spectacular. Assuming the photino (bino) is the NLSP, one has processes such as $e^+e^- \to \gamma\gamma + E\!\!\!/_t$ and $p\bar{b} \to e^+e^-\gamma\gamma + E\!\!\!/_t$, as indicated in Fig. 14.2, where$E\!\!\!/_t$ is the missing transverse energy.

Suggested reading

There are a number of good reviews of dynamical supersymmetry breaking, including those of Shadmi and Shirman (2000) and Terning (2003). The former includes catalogs of models and mechanisms. The recent interest in metastable supersymmetry breaking was launched by Intriligator *et al.* (2006). There is a large literature on gauge-mediated models

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and their phenomenology; a good review is provided by Giudice and Rattazzi (1999). The recent development of General Gauge Mediation is described in Meade *et al.* (2008). Models which achieve the full set of parameters are described in Buican *et al.* (2009) and Carpenter *et al.* (2009). A clear exposition of the origin of anomaly mediation is provided in Bagger *et al*. (2000), in Weinberg's text (1995), and in the more recent work of Dine and Seiberg (2007), Dine and Draper (2013), and DiPietro *et al.* (2014).

Exercises

(1) Check that the $SU(N)$ models, with an antisymmetric tensor and $N-4$ antifundamentals, have no flat directions and that they have a non-anomalous $U(1)$ symmetry.

(2) Verify Eq. (14.3) for the case of a $U(1)$ gauge theory with charged field ϕ^+ and ϕ^- introducing a Pauli–Villars regulator field.

(3) Check that Eqn. (14.5) is the most general expression that is consistent with symmetries, at least up to terms linear in m. Verify that there is no supersymmetric vacuum for this superpotential.

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