hep-th/yymmnnn
SCIPP-2005/03

Supersymmetry Breaking and Metastability

Michael Dine

Santa Cruz Institute for Particle Physics, Santa Cruz CA 95064

Abstract

Since *Supersymmetry and String Theory: Beyond the Standard Model* went to press, there have been a number of important developments in the subject of Dynamical Supersymmetry Breaking. These are reviewed here.

1 Dynamical Supersymmetry Breaking

Models with Dynamical Supersymmetry Breaking (DSB) known for some time. But until recently, dynamical supersymmetry breaking seemed an exceptional phenomenon, involving chiral gauge theories with a special structure. The resulting particle physics models suffered from various deficiencies:

- If DSB in a hidden sector in supergravity, gauginos light, and no amelioration of the flavor problems of ordinary supergravity theories.
- If DSB at lower scale, gauge mediation. But resulting models quite complicated. Most models have DSB scale at least 10^4 times the scale of weak interactions.

1.1 The ISS Model

Intriligator, Shih and Seiberg (ISS)[1] discovered susy breaking in a surprising context: a vectorlike theory, in which the Witten index was known to be non-zero. Their model is just SUSY QCD, with $N_f > N+1$, and *massive* quarks. Such theories were thought not to break supersymmetry because:

1. The Witten index, $\Delta = N$.
2. In the limit of large mass, this theory reduces to a pure, SUSY gauge theory, which is known not to break susy. For many questions, holomorphy of quantities such as the superpotential is enough to insure that there are supersymmetric minima.

Ironically, such gauge theories were originally proposed (with vanishing quark masses) as models of dynamical supersymmetry breaking; Witten's computation of the index was viewed as clinching the case that such theories did not break supersymmetry.

The breaking of supersymmetry is readily understood by considering the theory in its magnetic phase. In the infrared, the massless theory is described by an $SU(N_f - N)$ gauge theory with N_f q and $\bar{q}$ fields in the fundamental, and a set of "mesons", $\Phi_{f\bar{f}}$. The theory has superpotential:

$$W = \bar{q}\Phi q. \tag{1}$$

$\bar{q}$, q transform as $(\bar{N}_f, 1)$, $(1, \bar{N}_f)$ under the $SU(N_f) \times SU(N_f)$ flavor symmetry of the theory in the electric phase, without the mass term. Φ transforms as an $(\bar{N}_f, N_f)$. Under the $U(1)_R$, $\bar{q}, q$ carry charge N/N_f; Φ carries charge $2(1 - N/N_f)$.

1.1.1 Including the Mass Term

For small m, treat mass term as a perturbation. Transforms as $(\bar{N}_f, N_f)_{2(1-N/N_f)}$.

So in the magnetic theory,

$$\delta W = h \text{Tr } m \Lambda \Phi \tag{2}$$

has the same transformation numbers under the flavor symmetries as the underlying quark mass term.

1.1.2 SUSY Breaking in the Magnetic Theory

An important feature of the magnetic theory is that the kinetic terms for the dual quarks and the mesons are non-singular, and near the origin they can be taken to be canonical. It is a simple matter to check that the potential has a stable local minimum near the origin. There is a supersymmetric minimum away from the origin, which moves farther away as $m \to 0$.

E.g. $N = 2,\ N_f = 4$. Here the equation

$$\frac{\partial W}{\partial \Phi} = 0 \tag{3}$$

requires that $\bar{q}_f q_f'$ be proportional to the unit matrix. But q_f is a 2×4 matrix. By a symmetry transformation (and field redefinitions), this may be brought to the form:

$$\bar{q}_f q_f' = \begin{pmatrix} v^2 & 0 & 0 & 0 \\ 0 & v^2 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix}. \tag{4}$$

So one cannot satisfy the condition, and supersymmetry is broken.

1.1.3 What about the Witten Index?

Turn off m. Under the continuous R symmetry:

$$\Phi \rightarrow e^{2i\alpha(N_f - N)/N_f}\Phi \tag{5}$$

so expect non-perturbative superpotential:

$$W = \Lambda^{(-3N+N_f)/(N_f-N)} \mathrm{det}\Phi^{1/(N_f-N)}. \tag{6}$$

1.1.4 ISS As Basis of a Phenomenology

One loop computation: Potential for Φ has minimum at origin. For a range of parameters, the vacuum with broken susy is highly metastable.

In the presence of the mass term, the original theory has a discrete, Z_N R symmetry ($\alpha = 2\pi/N$). q, $\bar{q}$ are neutral under this symmetry. Φ transforms. The symmetry is unbroken in the magnetic phase. This R symmetry forbids a mass for gauginos.

A number of authors have constructed models in which the R symmetry is broken[2, 3, 4, 5, 6].

2 Retrofitting O'Raifeartaigh Models

The ISS models are part of a large class of models with broken supersymmetry in metastable minima. Any O'Raifeartaigh Model can be converted into a dynamical model Consider a model with chiral fields, A, Y and Z

$$W = \lambda Z(A^2 - \mu^2) + \lambda' Y A^2. \tag{7}$$

Generate μ dynamically. SU(N) pure gauge theory:

$$\int d^2\theta \frac{Z}{4M} W_\alpha^2. \tag{8}$$

Integrating out the gauge fields, leaves a superpotential:

$$W = \lambda Z A^2 + \frac{\Lambda^3 e^{-8\pi Z/b_0}}{M} + \lambda' Y A^2 \tag{9}$$

This structure is natural, in the sense that one can account for it with discrete symmetries.The gauge theory has a Z_N discrete symmetry, so if Z and Y are neutral, and A transforms like λ (with phase $e^{2\pi i/N}$, the only couplings of dimension three or less which are invariant are those above.

The original O'Raifeartaigh theory had a flat direction classically:

$$V = |F_Z|^2 \tag{10}$$

independent of Z. At one loop, Coleman-Weinberg calculation gives a minimum at $V = 0$.

$$W = \lambda Z A^2 + \frac{\Lambda^3 e^{-8\pi Z/b_0}}{M} + \lambda' Y A^2 \tag{11}$$

has a (supersymmetric) minimum at $Z \to \infty$ (runaway). Combined with the CW calculation, local minimum at origin, susy minimum at ∞.

2.1 Gravity Mediation

Already as models of gravity mediation, these theories are interesting. Couple to (super) gravity, along with fields of MSSM. Introduce constant in superpotential to tune cosmological constant. Then squarks and sleptons gain mass at tree level. No symmetry prevents coupling of Z to W_α^2 (unlike simplest conventional models of DSB). So gaugino masses also at tree level.

These models, while providing a dynamical explanation of the hierarchy, still have the standard difficulty of gravity mediation: They offer no insight into flavor problems.

2.2 Gauge Mediation

These are suitable models of gauge mediation. In the simple model we have described, the dynamical scale, μ, is a parameter, and so is the scale of susy breaking. But the low energy theory possesses an unbroken R symmetry. This is a feature of all models in this class (with chiral fields only coupled to a single gauge interaction). So we need to enlarge the class of models. A natural direction is the addition of gauge interactions.

2.3 O'Raifeartaigh Models with More Scales

In our simple model, the scale μ arose as $\mu^2 = \Lambda^3/M_p$. By coupling fields to higher dimension operators, we can generate other combinations. Another standard O'Raifeartaigh model is:

$$W = Z(A^2 - \mu^2) + M\ AB \tag{12}$$

Replace scales by:

$$Z(A^2 - \frac{W_\alpha^4}{M_p^4}) + AB\frac{W_\alpha^2}{M_p^2}. \tag{13}$$

Note $M \sim \mu \sim \Lambda^3/M_p^2$.

2.4 Models with Gauge Interactions

Adding a gauge symmetry changes the structure of the Coleman-Weinberg potential. A simple model with a $U(1)$ gauge interaction and charged chiral fields $Z^\pm$, $\phi^\pm$, and neutral field Z^0.

$$W = M(Z^+\phi^- + Z^-\phi^+) + \lambda Z^0(\phi^+\phi^- - \mu^2). \tag{14}$$

This breaks SUSY:

$$\phi^+ = \phi^- = v \qquad v^2 = \frac{\lambda^2\mu^2 - M^2}{\lambda^2} \tag{15}$$

(up to phases) while

$$F_{Z^+} = F_{Z^-} = Mv; \quad F_{Z^0} = \frac{M^2}{\lambda}. \tag{16}$$

There is a flat direction with

$$Z^\pm = -\frac{\lambda Z^0\phi^\pm}{M}. \tag{17}$$

Again, need to do Coleman-Weinberg calculation. For large Z, one can do the calculation easily using supergraphs. Work to second order in $F_{Z^\pm}$, F_{Z^0}. Study diagrams with two external Z fields. Result is:

$$V = \frac{1}{16\pi^2}\int d^4\theta(\lambda^2 Z^{0\dagger}Z^0 - 4g^2(Z^{+\dagger}Z^+ + Z^{-\dagger}Z^-))\ln(|Z|^2). \tag{18}$$

To obtain a sensible minimum, the coefficient of the log term must be positive.

For the full calculation of the potential, it is necessary to diagonalize mass matrix exactly, and use the Coleman-Weinberg formula:

$$V = \frac{1}{64\pi^2} \sum (-1)^F m^4 \ln(m^2). \tag{19}$$

The result of a straightforward, if somewhat tedious analysis (most sensibly done using Maple, Mathematica, or similar tools) is indicated in fig. zpotential.

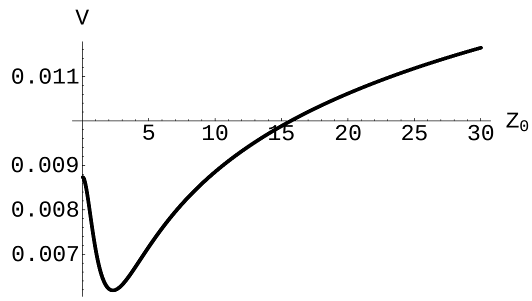

Figure 1: Z potential with $g = .4, \lambda = 1, M = 1, \mu = 1.5$.

One can now build a gauge-mediated model. The field Z^0 has non-zero scalar and auxiliary components. If coupled to the 5 and $\bar{5}$ of messengers.

$$\lambda' Z^0 \bar{M} M. \tag{20}$$

then for squarks, sleptons and gauginos we obtain the usual gauge-mediated spectrum. The scale is a parameter we are free to choose, and can vary wildly.

References

[1]

[2]

[3]

[4]

[5]

[6]

[7]

[8]