DYNAMICAL SUPERSYMMETRY BREAKING
IN SUPERGRAVITY THEORIES

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ABSTRACT

In rigid supersymmetry, generic models of dynamical supersymmetry breaking contain a light Goldstone boson, called the $R$ axion. We show that supergravity effects explicitly break the $R$ symmetry and give mass to the $R$ axion. For visible and renormalizable hidden sector models, the massive $R$ axion is free from astrophysical and cosmological problems. For nonrenormalizable hidden sector models, the $R$ axion suffers from cosmological difficulties similar to those of the moduli fields in string theory.

1. Introduction

Supersymmetry has long been viewed as an attractive candidate for physics beyond the standard model. There are many reasons for this, including the fact that supersymmetry stabilizes the gauge hierarchy against radiative corrections. In a supersymmetric theory, once the weak scale is fixed to be much smaller than the Planck scale, the hierarchy $M_W \ll M_P$ is preserved by radiative corrections.

In fact, the hierarchy can be preserved even if supersymmetry is broken. This follows from the fact that divergent radiative corrections are cut off by the scale of supersymmetry breaking. The hierarchy is preserved if the radiative corrections obey the following naturalness condition,

$$\delta M_W^2 \simeq \eta^2 M_S^2 \simeq M_W^2,$$

where $M_S$ is the scale of supersymmetry breaking and $\eta$ is an effective coupling that parametrizes the strength with which the supersymmetry breaking is communicated to the everyday world.

Although supersymmetry stabilizes the hierarchy, it does not necessarily explain the origin of the ratio

$$\frac{M_W}{M_P} \simeq 10^{-17}. \quad (2)$$

This motivates one to consider models of dynamical symmetry breaking, where supersymmetry (and electroweak symmetry) are dynamically broken. In such models all scales arise from $M_P$ through dimensional transmutation,

$$\Lambda \simeq \exp \left( \frac{-8\pi^2}{g^2} \right) M_P. \quad (3)$$
Supersymmetry is unbroken to all orders of perturbation theory, but nonperturbative
effects generate a ground state that breaks supersymmetry \[1, 2\].

Models with dynamical supersymmetry breaking offer a natural explanation for
the origin of the gauge hierarchy. Typically, such theories contain two sectors. One
contains the minimal supersymmetric standard model (or some simple extension).
The other dynamically breaks supersymmetry. The full theory can then be classified
by the way in which supersymmetry breaking is transmitted to the fields of the
standard model \[3\]:

- **Visible Sector Models.** In VS models, supersymmetry breaking is communicated
  by gauge interactions. Supersymmetry is broken by the dynamical theory at
  the scale \( \Lambda \), so \( M_S^2 \simeq \Lambda^2 \). The effective coupling \( \eta \) is of order \( (g/4\pi)^n \), where
  \( n \) counts the number of loops necessary to connect the two sectors. In such
  models, \( \Lambda \) is typically of order \( 10^5 \) GeV.

- **Renormalizable Hidden Sector Models.** In RHS models, supersymmetry break-
  ing is transmitted by gravitational interactions. Supersymmetry is still bro-
  ken by a dynamical theory at the scale \( \Lambda \), so \( M_S^2 \simeq \Lambda^2 \). Now, however,
  \( M_S^2 \simeq M_W M_P \), so \( \Lambda \simeq 10^{10} \) GeV. In these models, the effective coupling is
  much smaller, \( \eta^2 \simeq M_W/M_P \).

- **Nonrenormalizable Hidden Sector Models.** In NRHS models, the supersymmetry
  breaking is communicated by gravity, so \( M_S^2 \simeq M_W M_P \). The difference is that
  supersymmetry is not broken in the limit \( M_P \to \infty \). Typically, \( M_S^2 \simeq \Lambda^3/M_P \),
  which implies \( \Lambda \simeq 10^{13} \) GeV. NRHS models are not renormalizable in the sense
  that supersymmetry breaking relies on nonrenormalizable operators suppressed
  by powers of \( 1/M_P \). As in the previous case, \( \eta^2 \simeq M_W/M_P \).

In each of these cases, the dynamical supersymmetry breaking has important
consequences for the everyday world. These effects can be summarized in terms of a
spurion in an effective superspace lagrangian. Let us examine each type of model in
turn.

1.1. Visible Sector Models

In VS models, the supersymmetry breaking can be understood in terms of a \( \text{SU}(3) \times \text{SU}(2) \times \text{U}(1) \) chiral superfield \( S \). Below the scale \( \Lambda \), this superfield appears in the
effective lagrangian through terms of the form

\[
\frac{\eta^2}{\Lambda^2} \int d^4 \theta \ S^+ S \Phi^+ \Phi + \frac{\eta}{\Lambda} \int d^2 \theta \ SW^\alpha W_\alpha ,
\]

where \( \Phi \) and \( W^\alpha \) are matter and gauge superfields of the supersymmetric standard
model. The field \( S \) is a spurion, so it has no dynamics, just an expectation value
\( \langle S \rangle \simeq M_S^2 \theta \theta \). Therefore these terms give rise to soft supersymmetry-breaking masses of order \( \eta M_S \) for the squarks, sleptons and gauginos.

In VS models, the soft masses are, in principle, calculable. However, most models are quite complicated, with the soft masses arising at multi-loop order. (See, for example, \[4\]. For recent progress, see \[5\].) Therefore the effective coupling \( \eta \) tends to be small, perhaps of order \( 10^{-3} \), in which case \( M_S \simeq \Lambda \simeq 10^5 \text{ GeV} \).

1.2. Renormalizable Hidden Sector Models

In RHS models, supersymmetry breaking is communicated by the supergravity auxiliary fields. The superdeterminants of the superspace vielbeins play the role of the spurions,

\[
\int d^4 \theta ~ E \Phi^+ \Phi \\
\int d^2 \theta \ E W^\alpha W_\alpha ,
\]

where we use the notation of \[6\]. The normalizations of these terms are fixed by the normalizations of the superspace kinetic energies.

When supersymmetry is broken, the fields \( E \) and \( \mathcal{E} \) develop expectation values,

\[
\langle E \rangle \simeq 1 - \frac{M_S^2}{M_P} \theta \theta - \frac{M_S^2}{M_P} \bar{\theta} \bar{\theta} + \frac{M_S^4}{M_P^2} \theta \theta \bar{\theta} \bar{\theta}
\]

\[
\langle \mathcal{E} \rangle \simeq 1 - 3 \frac{M_S^2}{M_P} \theta \theta .
\]

The spurion \( E \) induces universal masses of order \( M_S^2/M_P \simeq M_W \) for the squarks and the sleptons. Therefore, in these models, \( \Lambda \simeq M_S \simeq 10^{10} \text{ GeV} \).

As in the VS case, RHS models have the advantage that they are, in principle, calculable. The problem is that \( \Phi \) gives a vanishing tree-level gaugino mass. (For two viewpoints on the phenomenology of light gluinos, see \[7\].) This follows from the fact that the supergravity auxiliary field that appears in \( \mathcal{E} \) also appears in the supersymmetric gauge field strength \( W_\alpha = (\bar{D} D - 8 R) e^{-V} D_\alpha e^V \). A careful analysis shows that it cancels between the two terms, rendering all gauginos massless. (For more discussion of RHS models, see \[8\].)

1.3. Nonrenormalizable Hidden Sector Models

The prototypical NRHS model is motivated by superstring theory, in which supersymmetry breaking is communicated by a dilaton field \( S \). (We ignore the additional complications from string moduli. See, for example, \[9\], and references therein.) The effective lagrangian contains terms of the form

\[
\int d^4 \theta ~ E \left\{ - M_P^2 \log \left( \frac{S + S^+}{M_P} \right) + \Phi^+ \Phi \right\}
\]
\[ \frac{1}{M_p} \int d^2 \theta \, \mathcal{E} S W^\alpha W_\alpha. \]  

(7)

Now, in a NRHS model, the gauginos are assumed to condense, \( \langle \lambda^\alpha \lambda_\alpha \rangle = \langle W^\alpha W_\alpha \rangle \) \( \simeq \Lambda^3 \). From (7) we see that this induces an expectation value of order \( \Lambda^3/M_p \) for the \( \theta \theta \)-component of \( S \). This, in turn, implies that supersymmetry is spontaneously broken.

The great advantage of this scenario is that it is motivated by a very general feature of string theory, the presence of a dilaton. It induces tree-level, weak-scale masses for the gauginos and the squarks and sleptons. The disadvantage is that the models are not calculable. Typically, there is no stable vacuum. At best there is a cosmological solution, in which the vacuum rolls off to the Planck scale, where it is assumed to be stabilized by unknown string effects. (See, however, [10].)

1.4. Common Features

At first glance, these three pictures differ considerably. However, for generic models – that is, models whose superpotentials contain all couplings allowed by symmetry – they share a common feature: the existence of an extra, unwanted, global, U(1) symmetry, called \( R \) symmetry [11]. When supersymmetry is broken, the \( R \) symmetry is also broken. This gives rise to a massless Goldstone boson, called the \( R \) axion. Such a particle is unacceptable for a variety of reasons. (See [12], and references therein.) The \( R \) axion must be eliminated, one way or another.

One way to eliminate the \( R \) axion is to construct models that are not generic. In supersymmetric theories, this can actually be natural [13], and much progress has been made in analyzing the properties of such theories [14]. Another escape is to add more structure so that the \( R \) axion becomes a pseudo-Goldstone boson [4]. This route leads to VS models that are quite complex.

In the rest of this talk, I will first explain why \( R \) symmetry is associated with generic models of dynamical supersymmetry breaking [11]. I will then show – for VS and RHS models – how supergravity effects eliminate the problems with the \( R \) axion [15]. I will conclude by illustrating how this works in the context of a specific model, the original 3 – 2 model of Affleck, Dine and Seiberg [2]. This model provides a nice test case because the effects of dynamical supersymmetry breaking can be calculated in a controlled weak-coupling expansion.

2. The Ubiquity of \( R \) Symmetry

To understand the role of \( R \) symmetry in dynamical supersymmetry breaking, let us suppose that we have a supersymmetric theory whose coupling grows strong at the scale \( \Lambda \). Let us consider the effective theory, valid below the scale \( \Lambda \). Let us also assume that the effective theory is supersymmetric, with the scale of supersymmetry breaking \( M_S \lesssim \Lambda \).

In this situation, the effective theory can be described by a Kähler potential \( K \) and superpotential \( W \). The superpotential is an analytic function of the fields \( z_i \), for
Supersymmetry is spontaneously broken if

$$\frac{\partial W}{\partial z_i} \neq 0$$

for some $i = 1,\ldots, n$.

In general, the system

$$\frac{\partial W}{\partial z_i} = 0$$

has a solution because it contains $n$ complex equations in $n$ complex unknowns. This implies that, for generic superpotentials, it is typically not possible to break supersymmetry.

The situation does not change if $W$ is invariant under a $d$-dimensional internal symmetry group. In this case the superpotential depends on $n-d$ complex variables, so (9) reduces to $n-d$ complex equations in $n-d$ unknowns. Generically, supersymmetry is not broken.

Now, however, let us assume that the effective theory has a spontaneously-broken continuous $R$ symmetry. Under an $R$ symmetry, the superpotential is not invariant, but has $R$-charge $-2$,

$$W \rightarrow e^{-2i\alpha} W .$$

Since the $R$ symmetry is spontaneously broken, we are free to label our fields so that the field $z_n$ has $R$-charge $q_n \neq 0$, with $\langle z_n \rangle \neq 0$. We can then write

$$W = z_n^{-2/q_n} F(x_j) ,$$

where $x_j = z_j/z_n$, for $j = 1,\ldots, n-1$. In terms of the variables $x_j$, the conditions (9) become

$$\frac{\partial F}{\partial x_j} = 0$$

$$F = 0 .$$

This system contains $n$ complex equations in $n-1$ unknowns. It does not generally have a solution, so supersymmetry is spontaneously broken.

This argument shows that – for generic models – spontaneously-broken $R$ symmetry is sufficient for dynamical supersymmetry breaking. As discussed above, the $R$ symmetry leads to a massless Goldstone boson – the $R$ axion. In what follows, we will show that supergravity effects explicitly break the $R$ symmetry and give mass to the $R$ axion.

### 3. The Supergravity Solution

In this section we shall see that supergravity provides a natural solution to the problem of the $R$ axion. To begin, we recall that in rigid supersymmetry, the scalar potential $V$ can be written in terms of the superpotential $W$ and the Kähler metric $K_{ij}$,

$$V = K^{ij*} \partial_i W \partial_j W^* ,$$

(13)
where, for simplicity, we ignore possible $D$-terms. Supersymmetry is spontaneously broken if the vacuum energy is positive, $\langle \partial_i W \rangle \neq 0$, for some value of $i = 1, \ldots, n$.

When supergravity is included, the scalar potential changes in the following way,

$$V = \exp \left( \frac{K}{M_P^2} \right) \left[ K^{ij} D_i W D_j W^* - \frac{3}{M_P^2} W^* W \right].$$  \hspace{1cm} (14)

In this expression, the covariant derivative of the superpotential is given by

$$D_i W = \partial_i W + \frac{K}{M_P} W.$$  \hspace{1cm} (15)

The condition for supersymmetry breaking is $\langle D_i W \rangle \neq 0$.

Comparing (13) with (14), we see that all of the supergravity corrections are suppressed by powers of $1/M_P$. The only term of consequence is a possible constant in the superpotential,

$$W \rightarrow W + c.$$  \hspace{1cm} (16)

In rigid supersymmetry, the constant has no effect and can safely be ignored. In local supersymmetry, however, the story is different. This is because the constant contributes to the vacuum energy. In fact, its role is to cancel the vacuum energy and ensure that the cosmological constant is zero. (In theories where all scales are less than $M_P$, this is the only way to cancel the vacuum energy.)

Given $\langle D_i W \rangle \simeq M_5^2$, the constant $c$ must be of order $M_5^2 M_P$ to cancel the vacuum energy. Because it grows with $M_P$, the constant is very important. It cancels the cosmological constant and generates the gravitino mass. It induces soft supersymmetry-breaking masses for the squarks and sleptons. But most importantly for this talk, it explicitly breaks the $R$ symmetry and gives rise to an explicit mass for the $R$ axion.

The mass of the $R$ axion is easy to find using a nonlinear realization. Under a field-dependent $R$ transformation, the superpotential $W$ transforms as follows,

$$W + c \rightarrow e^{2 i A / f_A} W + c,$$  \hspace{1cm} (17)

where $A$ is the axion field, and $f_A$ is the axion decay constant. The mass of the axion can be found by substituting (17) into (14) and expanding to second order in $A$. One finds

$$M_A^2 \simeq \frac{c}{f_A^2 M_P^2} \langle W \rangle,$$  \hspace{1cm} (18)

which implies

$$M_A^2 \simeq \frac{M_5^2 \Lambda^3}{f_A M_P}.$$  \hspace{1cm} (19)

For VS models, $f_A \simeq \Lambda \simeq M_S$, so $M_A^2 \simeq \Lambda^3 / M_P$. Astrophysical constraints based on stellar and supernova evolution require $M_A \gtrsim 10$ MeV [17], which implies $\Lambda \gtrsim 10^5$ GeV. Therefore VS models are safe provided $\Lambda$ is above about $10^5$ GeV, as is typically the case.
Table 1: The fields of the $3 - 2$ model.

| Particle | SU(3) $\times$ SU(2) | Hypercharge | R-charge |
|----------|----------------------|-------------|----------|
| $Q$      | (3, 2)               | 1/6         | 1        |
| $\bar{U}$ | (3, 1)              | $-2/3$      | 0        |
| $\bar{D}$ | (3, 1)             | $1/3$       | 0        |
| $L$      | (1, 2)               | $-1/2$      | $-3$     |

For RHS models, $f_A \simeq \Lambda \simeq M_S$, so $M_A^2 \simeq \Lambda^3/M_P$, as before. Now, however, $M_A^2 \simeq M_W M_P$, so $M_A^2 \simeq M_W M_S$, or $M_A \simeq 10^6$ GeV. Such an axion is too heavy to affect stellar dynamics. It decays relatively quickly so it is also cosmologically safe. (See [13], and references therein.) Therefore RHS models do not have problems with the $R$ axion.

Finally, for NRHS models, $f_A \simeq M_P$ and $M_A^2 \simeq \Lambda^3/M_P$. For such models, $M_A \simeq M_W$, about 100 GeV. Such a light, weakly coupled axion is cosmologically dangerous, but no more so than the light moduli fields that arise in string theory [17]. Presumably, the mechanism that cures the cosmological moduli problem also cures the cosmological problem with the $R$ axion. (For recent progress in this direction, see [18].)

4. The $3 - 2$ Model

The so-called $3 - 2$ model of Affleck, Dine and Seiberg provides a classic example in which dynamical supersymmetry breaking is realized in a controlled, weak-coupling expansion [2]. The model can serve as a prototypical RHS theory, or, suitably generalized, as the basis of a VS theory of dynamical supersymmetry breaking [3].

4.1. The Model

The $3 - 2$ model is based on two-flavor supersymmetric QCD, with a gauged SU(2)$_L$ flavor symmetry. To describe the model, let us denote the left- and right-handed quark superfields as $Q$ and $\bar{Q} = (\bar{U}, \bar{D})$. Under the SU(3) $\times$ SU(2) gauge symmetry, the quark superfields transform as shown in Table 1.

Note that the particle content of the model is similar to that of the minimal supersymmetric standard model, without the Higgs and the right-handed electron superfields. Apart from the gauge symmetries, the model also has two anomaly-free continuous global symmetries: hypercharge, U(1)$_Y$, and an $R$ symmetry, U(1)$_R$. The hypercharge and $R$-charge assignments are also listed in Table 1.

The Kähler potential of the model takes the usual form

$$K = Q^+ Q + \bar{Q} \bar{Q}^+ + L^+ L.$$  \hspace{1cm} (20)

In (20) the SU(2) and SU(3) gauge superfields are not written, but are assumed to
Table 2: The fields of the effective theory.

| Particle | Hypercharge | $R$-charge |
|----------|-------------|------------|
| $X_1$    | 0           | $-2$       |
| $X_2$    | $-1$        | $-2$       |
| $X_3$    | 0           | 2          |

be coupled in the usual way [6].

In the absence of a superpotential, the scalar potential vanishes for a number of flat directions in field space. Therefore the ground state is undetermined at the classical level [2]. The equations that determine the flat directions are

$$Q^+ m Q_\ell - Q^m Q^\ell_+ = 0$$  \hspace{1cm} (21)

for the SU(3) $D$-terms, and

$$Q_{\alpha}^+ Q^\beta + L_{\alpha}^+ L^\beta = \frac{1}{2} \delta_{\alpha}^\beta (Q^+ Q + L^+ L)$$  \hspace{1cm} (22)

for the SU(2) $D$-terms. Up to local symmetries, the solutions to these equations are parametrized by six real variables, also called moduli. The moduli parametrize inequivalent, supersymmetry-preserving vacua. The variations along the moduli directions correspond to six real, massless scalar fields.

Let us now consider the theory expanded around a solution of (21), (22), such that the vacuum expectation values $v$ obey

$$v \gg \Lambda,$$  \hspace{1cm} (23)

where

$$\Lambda = v \exp \left( -\frac{8\pi^2}{g(v)^2 b_0} \right).$$  \hspace{1cm} (24)

In this expression, $\Lambda$ is the scale where the SU(3) gauge coupling $g$ becomes strong, and $b_0$ is the one-loop coefficient of the SU(3) beta function. For $v \gg \Lambda$, the theory is in the weak-coupling regime. The SU(3) and SU(2) gauge symmetries are completely broken so the vector supermultiplets are massive. Supersymmetry is unbroken, so 11 out of the 14 matter chiral superfields are absorbed by the vector superfields. The six real moduli are contained in three massless chiral superfields.

At energies below the scale $\Lambda$, the low-energy effective theory can be described in terms of three gauge-invariant chiral superfields,

$$X_1 = Q D L,$$
$$X_2 = Q U L,$$
$$X_3 = \det Q_\alpha Q^\beta,$$  \hspace{1cm} (25)
whose scalar components parametrize the six massless moduli. The quantum numbers of these fields are listed in Table 2.

Let us now discuss the superpotential of the effective theory. The vacuum preserves the global hypercharge symmetry, so we take the superpotential to preserve it as well,

$$ W = \lambda X_1 + 2 \frac{\Lambda^7}{X_3}. $$ (26)

The first term is the renormalizable superpotential that we assume to be present in the classical theory. The second is nonrenormalizable, and can be shown to be generated by nonperturbative effects. Its coefficient can be calculated in a weak coupling expansion around a constrained instanton vacuum [2]. Note that this superpotential also respects U(1)$_R$. The $R$ symmetry is “accidental” in the sense that it is a direct result of hypercharge conservation in the effective theory.

In the presence of the superpotential, the scalar potential is no longer flat. Indeed, when $\lambda = 0$, the scalar potential does not have a minimum and the theory does not have a well-defined ground state [2]. For the case when

$$ \lambda \ll g_2 \ll g_3 \ll 1, $$ (27)

the potential has a minimum at finite values for the fields, of order

$$ v \simeq \frac{\Lambda}{\lambda^{1/7}}, $$ (28)

as shown in Figure 1. This value is such that the weak coupling assumption (23) is self-consistent, so the theory can be analyzed perturbatively.

As we will see, the vacuum energy is positive and supersymmetry is spontaneously broken, with $M_S \simeq \lambda^{5/14} \Lambda$. The moduli are massive, with masses of order $\lambda^{6/7} \Lambda$. The hierarchy of scales is summarized in Figure 2.
4.2. The Low-Energy Sigma Model and its Spectrum

In this section we will study the effective field theory below the scale $\Lambda$. We will find the spectrum of all particles lighter than this scale.

In the limit (27), the Kähler potential of the effective theory is given by the projection of (20) onto the moduli fields $X_1, X_2$ and $X_3$. (See also [19].) Using the notation of ref. [2], the resulting Kähler potential can be written

$$K = 24 \frac{A + Bx}{x^2} ,$$

where

$$A = \frac{1}{2} (X_1^+ X_1 = X_2^+ X_2)$$

$$B = \frac{1}{3} \sqrt{X_3^+ X_3} ,$$

and

$$x = 4 \sqrt{B} \cos \left( \frac{1}{3} \text{Arccos} \frac{A}{B^{3/2}} \right) .$$

The equations that determine $x$ as a function of the light superfields have several solutions. Equation (31) is the only one that leads to a positive definite Kähler metric at the minimum.
The low-energy theory is therefore described by a sigma model with Kähler potential $K$ and superpotential (26)

$$W = \lambda X_1 + 2 \frac{\Lambda^7}{X_3} .$$  \hfill (32)

To find the ground state, one must minimize the scalar potential. After a numerical analysis, one finds the following values for the expectation values of the moduli fields:

$$X_1 = 0.50 \frac{\Lambda^3}{\lambda^{3/7}}$$

$$X_2 = 0$$

$$X_3 = 2.58 \frac{\Lambda^4}{\lambda^{4/7}} .$$  \hfill (33)

At the minimum, the vacuum energy density is

$$M_S^4 = 3.59 \frac{\lambda^{10/7}}{\Lambda} .$$  \hfill (34)

The scalar mass matrix is found by expanding the potential about its minimum. It is given by

$$M_{ab}^2 = \langle V_{ab} \rangle ,$$  \hfill (35)

where $V$ is the scalar potential, and $a, b = 1, \ldots, 6$ label the six light real fields. One discovers three real scalar fields of masses 3.88, 2.83 and 2.04 (in units of $\lambda^{6/7} \Lambda$), one complex scalar of mass 1.35 (in the same units), and a massless $R$ axion [15].

The fermion mass matrix is [16]

$$M_{ij} = \langle W_{ij} - K_{k\ell}^{-1} K_{ij\ell} W_k \rangle ,$$  \hfill (36)

where $i, j = 1, \ldots, 3$ label the three light fermions. One finds a massless goldstino, a massless fermion of hypercharge one, and a fermion of mass 3.19 $\lambda^{6/7} \Lambda$ [15].

4.3. Supergravity Couplings and $R$ Axion Mass

In this section we couple the $3-2$ model to supergravity and compute the supergravity contribution to the $R$ axion mass. The supergravity coupling is straightforward and can be done using the effective theory of the previous section.

As discussed above, in any supergravity theory where all scales are smaller than $M_P$, the vacuum energy is canceled by adding a constant to the superpotential,

$$c = \frac{1}{\sqrt{3}} M_S^2 M_P .$$  \hfill (37)

In this class of models, all soft breaking terms are induced by $c$. For the $3-2$ model, gravitino mass is

$$M_{3/2} = \frac{c}{M_P^2} + \frac{1}{\sqrt{3}} \frac{M_S^2}{M_P} = 1.09 \frac{\lambda^{5/7}}{M_P} .$$  \hfill (38)
For our purposes, the most important consequence of the constant $c$ is the fact that it explicitly breaks the $R$ symmetry. In particular, it induces $R$-symmetry-breaking terms in the supergravity-coupled scalar potential, including

$$V' = \frac{1}{\sqrt{3}} \frac{M_S^2}{M_P} \left( W_i K_{ij}^{-1} K_{j*} - 3 W \right) + \text{h.c.} + \ldots,$$

(39)

where $K$ and $W$ are the Kähler potential (29) and superpotential (32) of the effective theory, and the dots denote terms suppressed by additional powers of $M_P$.

As discussed above, these terms give mass to the $R$ axion. The axion coupling constant turns out to be

$$f_A = 2.18 \frac{\Lambda}{\Lambda^{1/7}} = 1.58 \frac{M_S}{\sqrt{\lambda}}.$$

(40)

The axion mass is then [13]

$$M_A^2 = 10.0 \frac{\Lambda^{11/7}}{M_P} \frac{\Lambda^3}{M_P} = 6.58 \sqrt{\lambda} \frac{M_3}{M_P} \frac{M_S}{M_P}.$$

(41)

This formula is in agreement with our previous results for VS and RHS models.

5. Conclusions

In this talk we have seen that models of dynamical supersymmetry breaking offer an attractive explanation for the ratio $M_W/M_P \simeq 10^{-17}$. We have seen why it is difficult to construct such models, and demonstrated why many candidate models contain a potential $R$ axion.

In VS models with a supersymmetry breaking scale greater than about $10^5$ GeV, the axion is sufficiently heavy to evade astrophysical constraints. In RHS models, the axion mass is quite large, of order $10^6$ GeV, so the axion is astrophysically and cosmologically safe. In NRHS models, the axion mass is of order 100 GeV. Such a light, weakly-coupled axion can lead to cosmological difficulties of the sort already present for the moduli fields of string theory.

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