F-term SUSY Breaking and Moduli

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We discuss the coupling of heavy moduli fields to light fields when the dynamics of the latter, absent such couplings, yields metastable vacua. We show that the survival of the vacuum structure of the local model depends nontrivially on the cross-couplings of the two sectors. In particular we find that for "local" models (such as those realized by D-branes in type II string theories) with metastable vacua breaking supersymmetry via F-terms, cross-coupling of the two sectors at an intermediate scale can push the metastable vacuum outside of the regime of the effective field theory. We parametrize the region in which the metastable vacua are safe. We the show that sufficiently small cross-couplings can be made natural. Finally, we briefly discuss the role of moduli in stringy realizations of "retrofitted" SUSY-breaking sectors.
1. Introduction

Supersymmetric field theories with dynamical or spontaneous SUSY-breaking via F-terms generically have an effective field theory description as an O’Raifeartaigh-like model. Such models consist of scalar fields which spontaneously break supersymmetry due to a "rank condition" on the superpotential. The classic example is the following superpotential for three chiral scalars:

$$W(X, A, B) = hXA^2 + mAB + \mu X$$

(1.1)

If $y = h\mu/m^2 < 1$, this has a single SUSY-breaking vacuum at $A = B = 0$, $F_X = \mu$ (see for example [1,2] for a thorough discussion of this model). For any value of $y$, $X$ is a flat direction at tree level and is stabilized at the $U(1)_R$-preserving point $X = 0$ by the one-loop Coleman-Weinberg potential. One may also deform this model by a small $R$-breaking term, $\delta W = \frac{1}{2}\epsilon X^2$, which leads to a SUSY-preserving vacuum at large $X = -\mu_0/\epsilon$ and leads to a metastable SUSY-breaking vacuum at $X \sim \epsilon f(h, m, \mu)$.

In string theory, the couplings $m, \mu, h$ generally depend on moduli such as parameters of the metric of the compactification manifold. For example, in the local orientifold model of [3], $X$ is an open string field, and the analog of $\mu$ is generated by D-brane instantons which depend on a Kähler class of the orientifold. This Kähler class will become a dynamical modulus in a compact model. We might expect that if the geometric moduli are stabilized at a high scale, they (a) do not interfere significantly with the SUSY-breaking dynamics, and (b) do not acquire significant F-terms. One could then engineer a scenario in which the SUSY breaking in the theory (1.1) is communicated to a supersymmetric extension of the Standard Model via gravity or gauge or anomaly mediation.

These conditions would seem to be a natural consequence of Wilsonian decoupling, but one must take a little care. To see this, consider a quadratic theory with two scalar fields and the following mass matrix:

$$\mathcal{M}^2 = \begin{pmatrix} M^2 & \gamma \\ \gamma & m_{light}^2 \end{pmatrix}$$

(1.2)

Here $M \gg m_{light}$, so it appears that we have a heavy "modulus" and a light field, and we can integrate the modulus out by setting it to zero. However, if the "cross-coupling" term $\gamma$ is larger than the intermediate scale $\sqrt{Mm_{light}}$, the system is tachyonic. We will discuss further examples in §2. Note that in this model, the tachyon can be removed by imposing a $\mathbb{Z}_2$ symmetry which forces $\gamma = 0$. 

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In §3 we will discuss the case of light fields in globally supersymmetric models whose dynamics, when decoupled from moduli, yield metastable vacua which break SUSY through F-terms. We will show that when the simplest class of models is coupled to moduli, there is a lower bound on the cross-coupling beyond which the vacuum is pushed out of the range of effective field theory. In §4 we will discuss natural values of the cross-coupling terms and show that for ”geometric” moduli which have large ranges, the cross-couplings can be naturally suppressed.

We will conclude by reviewing another possibility for generating a sector such as (1.1) [9-11], in which the dimensionful couplings are generated by nonperturbative gauge dynamics. In string theory, X will typically be a geometric modulus itself; this modulus is stabilized by a one-loop mass, and acquires an F-term. Such models can be realizations of ”moduli domination” [12-1].

2. Nonsupersymmetric examples

Cross-couplings between ”light” and ”heavy” fields can change the vacuum structure of a theory if those couplings are set by an intermediate mass scale. Let us consider examples with two real scalar fields, a ”light” field $x$ and a heavy ”modulus” $\phi$. We will consider quadratic kinetic terms, and will assume that we have included all quantum corrections in our effective action (although we will, for computational ease, write fairly nongeneric potentials).

First, consider the potential

$$V = \frac{1}{4} \lambda x^4 + \frac{1}{2} m^2 x^2 + \gamma x \phi + \frac{1}{2} M^2 \phi^2$$

(2.1)

where $M \gg m$. If we integrate out $\phi$, we find the modified potential

$$V = \lambda x^4 + \frac{1}{2} (m^2 - \frac{\gamma^2}{M^2}) x^2.$$  

(2.2)

Assume $m^2 > 0$, so that if we naively set $\phi = 0$, the potential (2.1) would have a single vacuum at the origin $x = 0$. If $\gamma$ is larger than the intermediate scale $mM$, the actual

1 There have been a number of general discussions of when SUSY-breaking vacua exist in local and globally supersymmetric models – see for example [14-1]. There have also been some general discussions of subtleties of integrating heavy fields in local and global SUSY models in [5-8]. We are concerned with the particular issue of the effects of coupling light fields with apparent SUSY-breaking dynamics to heavy moduli.
vacuum structure will change, with $x = 0$ becoming an unstable local maximum of the potential.

A more generic example begins with the following relatively generic potential for the light field $x$:

$$V = \frac{1}{4}\lambda x^4 + \frac{1}{3}g x^3 + \frac{1}{2}m x^2$$

(2.3)

where $\lambda > 0$. This has extrema at

$$x = 0, x_- = -\frac{g}{2} \pm \sqrt{\frac{g^2}{4\lambda^2} - \frac{m^2}{\lambda}}$$

(2.4)

So long as $m^2 \leq g^2/(4\lambda)$, there are three extrema, two local minima and a local maximum. For example, if $m^2 > 0$, the minima are at $x_-$ and 0, with $x_+$ the local maximum: the metastable local minimum at $x_-$ coalesces with the maximum at $x = 0$ and disappears as we dial $m^2 > g^2/4\lambda$.

Now, let us couple (2.4) to a "heavy" modulus $\phi$. We will assume that the marginal and relevant couplings between $\phi$ and $x$ are linear in $\phi$ (this is not generic – we are making this assumption for ease of illustration, though one might make the higher-order terms in $\phi$ small using the considerations of §4):

$$V = \frac{1}{4}\lambda x^4 + \frac{1}{3}(g + g_1 \phi) x^3 + \frac{1}{2}(m^2 + m_1 \phi)x^2 + \gamma \phi x + \frac{1}{2}M^2 \phi^2$$

(2.5)

In this case, we can integrate out $\phi$ exactly to find that:

$$V = -\frac{g_1^2}{18M^2} x^6 - \frac{g_1 m_1}{6M^2} x^5 + \frac{1}{4}\lambda_{eff} x^4 + \frac{1}{3}g_{eff} x^3 + \frac{1}{2}m_{eff} x^2$$

(2.6)

where

$$\lambda_{eff} = \lambda - \frac{4g_1 \gamma}{3M^2} - \frac{m_1^2}{M^2}$$

$$g_{eff} = g - \frac{3m_1 \gamma}{2M^2}$$

(2.7)

$$m_{eff} = m^2 - \frac{\gamma^2}{M^2}$$

The $x^6$ and $x^5$ terms will generally be important compared to the lower-order terms only when $x$ is of order $M$ or the initial couplings $\lambda, g, m$ are extremely small.

Concentrating on the quartic and lower effective interactions, it becomes clear that while one may change the sign of the quadratic term by choosing $\gamma$ at the intermediate scale $mM$, even then it will be difficult to change the sign of the cubic and quartic terms, or to achieve $m_{eff}^2 \geq g_{eff}^2/(4\lambda_{eff})$. For these latter changes, the massive coupling $m_1$ must
be of order $Mg/m$ or the dimensionless coupling $g_1$ is extremely large, of of order $\lambda M/m$). Nonetheless, the sign change of the quadratic term will have the effect of rendering the $x = 0$ vacuum unstable and the vacua at $x_{\pm}$ stable. Inspired this, we will now discuss supersymmetric vacua with metastable SUSY-breaking states, and ask when the coupling to moduli can change the vacuum structure.

3. Supersymmetric examples

In this section we will consider globally supersymmetric $N = 1$ theories which break supersymmetry spontaneously through F-terms before coupling to moduli. These models can be taken to represent the dynamics of open strings on D-branes which are localized in the internal geometry of a type II compactification; of course, there could be other representations. Before coupling to moduli, the SUSY-breaking vacua of these models may be completely stable or metastable; generically, the latter occurs when one perturbs the theory by operators, such as the superpotential term $\epsilon X^2$ added to (1.1), which break the $U(1)_R$ symmetry of the theory [1,2,17]. We will focus on a particular form of the cross-coupling (arising from letting the SUSY-breaking coupling in the superpotential depend on a heavy modulus), and find that this coupling can change the vacuum structure. First, either the cross-coupling or the moduli mass term will break the $U(1)_R$ symmetry of the SUSY-breaking sector, so that even if the model (1.1) has a SUSY-breaking global minimum, the full theory will have supersymmetric minima. Furthermore, the cross-couplings can push the SUSY-breaking metastable state out of the range of the effective field theory. We will discuss this in detail in the context of a Kähler-stabilized Polonyi model, and find that the physics of the SUSY-breaking sector can be significantly changed when the moduli mass $M$ is much smaller than the UV mass scale $\Lambda$ parametrizing the suppression of higher-order terms in the Kähler potential. We will briefly discuss the explicit case of the O’Raifeartaigh model as well (which at sufficiently low energies and for a range of parameters, is well-described by a Kähler-stabilized Polonyi model).

\footnote{2 The recent work [18] includes the effects of heavy moduli in generating a stable potential for the light fields. This work does not include the cross-couplings discussed in this section; however it is quite possible that they can be suppressed using the considerations of §4.}

\footnote{3 Recent results on the stability of SUSY-breaking vacua [19] concentrate on tree-level physics with a minimal Kähler potential, so that this falls outside of the considerations of that work.}
3.1. Polonyi model

We will consider a theory with Kähler potential

$$K = |X|^2 + |Z|^2 - \frac{c}{\Lambda^2}|X|^4$$

(3.1)

where $\Lambda$ is a cutoff scale such as the Planck scale, and $c$ is a dimensionless number, and superpotential:

$$W = (\mu_0 + \mu_1 Z)X - \frac{1}{2}\epsilon X^2 + \frac{1}{2} M Z^2$$

(3.2)

This is mean to model the dependence of $\mu X$ on moduli.

If we decouple the two fields by setting $\mu_1 = 0$, $X$ has a SUSY vacuum at $X = \mu_0/\epsilon$. So long as

$$\epsilon < \frac{\sqrt{2c\mu_0}}{\Lambda},$$

(3.3)

then there is a metastable SUSY-breaking vacuum at $X = \frac{\epsilon\Lambda^2}{2\mu_0 c}$, and

$$F_x^* = \mu_0 - \frac{\epsilon^2\Lambda^2}{2c\mu_0}$$

(3.4)

The second term remains smaller than the first if (3.3) is satisfied. For larger values of $\epsilon$, the vacuum would be at $X > \frac{\Lambda}{\sqrt{2c}}$: the kinetic term for $X$ that arises from (3.1) alone flips sign, and higher-order terms in the Kähler potential must be taken into account. This region is outside of the domain of our effective theory.

If we include the coupling to $Z$, we can integrate out $Z$ as follows \[7\]. The equations of motion for the components of $Z$ are:

$$F_z^* = \partial_z W = \mu_1 x + M z$$

$$-\partial^2 z = \mu_1 F_x + MF_z$$

(3.5)

If we assume that the fields are constant, we find that the potential is:

$$U(x) = \left| \mu_0 - \left( \frac{\mu_1^2}{M} + \epsilon \right) x \right|^2$$

$$1 + \frac{\mu_1^2}{M^2} - \frac{2c}{\Lambda^2}|x|^2$$

(3.6)

where $x$ is the scalar component of $X$. We will assume that $c, \mu_0, \mu_1$ and $\epsilon$ are all real.

\[4\] In this case, solving $\partial Z W = 0$ for $Z$ yields the same answer. In general, as pointed out in \[7\], this does not give the full potential, though the difference will be suppressed by powers of $M$. 

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The potential \((3.6)\) is singular at:

\[
|x_{\text{sing}}| = \frac{\Lambda}{\sqrt{2c}} \left( 1 + \frac{|\mu_1|^2}{M^2} \right)
\]  
\((3.7)\)

If all of the parameters are real, as we supposed above, the solutions to \(\partial_x U = 0\) are real. There is a supersymmetry-preserving solution at:

\[
x_{\text{susy}} = \frac{\mu_0}{\mu_1^2/M^2 + \epsilon}
\]  
\((3.8)\)

and SUSY-breaking solutions to \(\partial_X U = 0\) at \(x = \infty\) and at

\[
x_{\text{sb}} = \frac{\Lambda^2}{2c\mu_0} \left( 1 + \frac{\mu_1^2}{M^2} \right) \left( \epsilon + \frac{\mu_1^2}{M} \right)
\]  
\((3.9)\)

The solution to \(\partial_x U = 0\) between the origin and \(x_{\text{sing}}\) will be stable. The dynamics at larger field values is out of range of our effective theory; higher-order corrections in the Kähler potential and superpotential will become important.

\(X = x_{\text{sb}}\) will be the stable vacuum between the origin and \(x_{\text{sing}}\) if:

\[
\epsilon M + \mu_1^2 < \sqrt{2c} \frac{M}{\Lambda} \mu_0 \left( 1 + \frac{\mu_1^2}{M^2} \right)^{1/2}
\]  
\((3.10)\)

Let us assume that \(\mu_0^{1/2}, \mu_1, \epsilon \ll M, \Lambda\). If \(\epsilon\) is such that the SUSY-breaking vacuum would be stable if \(\mu_1 = 0\), there is still a condition on \(\mu_1\): if \(M/\Lambda \ll 1\), \(\mu_1\) must be considerably smaller than the mass scale set by \(\mu_0\), lest the coupling to \(Z\) pushes the SUSY-breaking vacuum out of the regime of our effective theory. In particular, since \(\sqrt{2c\mu_0}/\Lambda \sim m_x\) is the mass of \(X\) before coupling to \(Z\), this means that the SUSY-breaking vacuum is under control if \(\mu_1^2 < m_x M\).

We can also compute the corrections to the F-terms that arise from coupling \(X\) and \(Z\). Using \((3.5)\) and

\[
F_x^* = \mu_0 + \mu_1 Z - \epsilon X
\]  
\((3.11)\)

we find after a few lines of algebra that:

\[
F_x^* = \frac{\mu_0}{1 + \frac{\mu_1^2}{M^2}} - \frac{\Lambda^2}{M^2} \left( \epsilon + \frac{\mu_1^2}{M} \right)^2 \frac{1}{2c\mu_0}
\]
\((3.12)\)

\[
F_z^* = \frac{\mu_1}{M} F_x^*
\]
So long as (3.10) is satisfied, we can show the second term in the expression for $F^*_x$ is smaller than the first, and in particular the effect of coupling to $Z$ is small. Furthermore, so long as $M$ is larger than the parameters $\mu_0^{1/2}, \mu_1$ of the local model, $F_Z$ remains much smaller than $F_x$. On the other hand, if $M \sim \Lambda$, and the inequality (3.10) is close to being saturated, the F-term for $X$ is reduced substantially from the result $F^*_x = \mu_0$ of the local model.

We have focused on a particular form of the cross-coupling in globally supersymmetric models. It would be interesting to study cross-couplings in the Kähler potential itself. It would also be interesting to study this question in the context of supergravity – for example, even if the Kähler potential and superpotential factorize, supergravity effects will typically induce cross-couplings; furthermore, in may interesting examples, the modulus $Z$ will be stabilized by supergravity effects. For the theory above, supergravity will not induce any quadratic cross-couplings between $X$ and $Z$, and we expect the effects to be small if $\Lambda \ll m_{pl}$.

3.2. The O’Raifeartaigh model

Let us consider the O’Raifeartaigh model coupled to a modulus $z$:

$$W = hXA^2 + mAB - (\mu_0 + \mu_1 Z)X + \frac{1}{2}MZ^2$$

(3.13)

with canonical kinetic terms for $X, Z$. In principle we should let $m, h$ also depend on moduli; we avoid this for simplicity’s sake. At energies below $m$, we can integrate out $A, B$ and $Z$. For $h\mu_0/m^2 \ll 1$, when we are studying a SUSY-breaking vacuum close to $X = 0$, the ”Coleman-Weinberg” potential which arises from integrating out $A, B$ can be expressed as a correction to the Kähler potential. (See for example [1,2] for a thorough discussion of the Coleman-Weinberg potential in this model.)

Since $Z$ only appears quadratically in the superpotential it can be integrated out classically. Let us first integrate out $A, B$ (which we can do first because these fields do not couple to $Z$.) If $X$ remains close to the origin, the result is essentially a Polonyi model of the type discussed above (there is also an order $h^2$ shift in the coefficient of $X^2$ in (3.1)); the $|X|^4$ term will be of the form (3.1), with $2c = \frac{|h|^4}{64\pi^2}$ and $\Lambda^2 = m^2$. If $m, \mu_0^{1/2}, \mu_1 \ll M$, then the SUSY-breaking vacuum in the O’Raifeartaigh model is safely inside the realm of the original effective field theory.
4. Stability and naturalness

When are conditions such as (3.10), or $\gamma < mM$ in (2.1), satisfied? One possibility arises if, as with axions, the range of the moduli is large while the potential is generated at a somewhat lower scale. For example, geometric moduli have a kinetic term of the form

$$S = m^2_{pl,4} \int d^4x G_{ab}(\phi) \partial \phi^a \partial \phi^b$$

(4.1)

where $\phi$ are dimensionless, since they parametrize the metric of the compactification manifold. In many models such as type II flux compactifications or heterotic M-theory compactifications on a Calabi-Yau times a large interval, potentials then arise either from fluxes from D-brane instantons, or from gauge instantons, and have the form

$$V = M^4 v(\phi)$$

(4.2)

where $M \ll m_{pl,4}$. When we rescale $\phi^a \rightarrow z^a = m_{pl,4} \phi^a$, where $z^a$ are dimension-1 4d scalars with canonical kinetic terms, we find that

$$V = M^4 v \left( \frac{z}{m_{pl,4}} \right)$$

(4.3)

Thus interactions are suppressed by powers of $1/m_{pl,4}$.

In many of these examples, this suppression is the result of an $\mathcal{N} = 2$ supersymmetry (which would forbid moduli masses) that is broken by fluxes at a scale lower than the Kaluza-Klein scale \[23-26\], or broken by D-branes and orientifolds which are local on the internal manifold and so are expected to induce $\mathcal{N} = 1$ masses suppressed by volume factors. Note that if we are interested in the effects of Planck-suppressed couplings, we will also have to take into account additional terms in the potential which arise from supergravity corrections.

Similarly, assume that supersymmetry-breaking dynamics arise from ”local models” of open string fields trapped on D-branes (or in the case of heterotic M-theory, from the $E_8$ ”walls”). Tree-level couplings to dimension-$\Delta$ operators $\mathcal{O}$ in the open string sector take the form:

$$\delta L = \frac{m_{s}^{\Delta-4}}{g_s} f(\phi) \mathcal{O} = \frac{m_{s}^{\Delta-4}}{g_s} f \left( \frac{z}{m_{pl,4}} \right) \mathcal{O} ,$$

(4.4)

\[ See [20] for a discussion of this phenomenon in type II flux compactifications, and [21,22] for a similar argument in the context of heterotic M-theory.
where $m_s$ is the string scale (or in heterotic M-theory, the 11d Planck scale) which is generally less than the 4d Planck scale. Instanton-generated couplings such as those in \[3\] will have the form:

$$
\delta L = \frac{M_{np}^2}{g_s} f(\phi) \mathcal{O} = \frac{M_{np}^2}{g_s} f \left( \frac{z}{m_{pl,A}} \right) \mathcal{O}.
$$

(4.5)

where $M_{np}$ is nonperturbatively small compared to $m_s, m_{pl,A}$.

In particular, consider the case (4.3), and expand the theory in small fluctuations about the appropriate minimum of (4.3). The SUSY-breaking parameter will be $\mu_0 \sim M^2$, while $\mu_1 \sim \frac{M}{m_{pl,A}}$. So long as $M_{np}/m_{pl,A} \ll M/\Lambda$ (where $M$ is the moduli mass and $\Lambda$ the mass scale in (3.1)), (3.10) should be easily satisfied. (Again, if we wished to consider the small effects of these couplings, we should also consider couplings induced by supergravity effects.)

On the other hand, if we have a string theory model in which (4.3) is determined by (for example) some fluxes, then we can think of $\mu_0, \mu_1, c, \epsilon, \Lambda$ and so on in (3.10) as depending on the value of these fluxes through the value $Z = z_0$ about which one expands to derive the effective theory in (3.1),(3.2). If we replace $Z$ with $\delta Z = Z - z_0$ in (3.1),(3.2), then (3.10) carves out a region of the space of fluxes for which the effective field theory of $X, \delta Z$ describes SUSY-breaking dynamics.

Note that a problem potentially arises if the moduli mass is generated by the same instantons as the coupling of the modulus to the SUSY-breaking dynamics. In this case, $M \sim M_{np}^2/m_{pl,A} \sim \mu_1$, and (3.10) is at best marginally satisfied. Furthermore, according to (3.12), the F-term for the modulus would be of the same order as the F-terms for the "local model" of SUSY-breaking. This is closer to moduli-domination scenarios as described in [12-16]. In type II theories this could be avoided by burying the D-branes responsible for SUSY-breaking down a warped throat, while ensuring that there are also D-instantons supported away from the warped region which give the moduli a mass. In addition, we could forbid the leading-order couplings with a discrete symmetry; terms which are higher order in $\phi$ will be suppressed by additional powers of $m_{pl,A}$.

5. Conclusions

In the last two sections we have discussed cases where the SUSY-breaking dynamics arises from fields distinct from the geometric moduli. These might arise if the SUSY-breaking dynamics is "local": for example, if it is localized on (intersecting) D-branes in type II string theory, or on an "$E_8$ wall" in heterotic M-theory.
One can also imagine scenarios in which the moduli participate more directly. One scenario for designing naturally small values of \( m, \mu \) in the theory (1.1) is to couple the fields \( X, A, B \) to non-abelian gauge fields \([9-11]\), sometimes known as "retrofitting". Gaugino condensation introduces the small mass scales in such models. For example, consider the case:

\[
W = hX A^2 + \frac{a}{M_s} X W_\alpha W^\alpha + \frac{b}{M_s^2} AB \tilde{W}_\alpha \tilde{W}^\alpha
\]  

(5.1)

where \( W, \tilde{W} \) are the chiral Fermi superfields for two \( SU(2) \) gauge groups, and \( M_s \) is the cutoff scale of the theory, typically the Planck scale. Typically \( X, A, B \) are moduli themselves. For example, if \( W, \tilde{W} \) are open string gauge groups in type IIB string theory, the gauge couplings will depend on the Kähler moduli of the theory.

Gaugino condensation will lead to the following low-energy superpotential:

\[
W = hX A^2 + \Lambda^3 e^{\frac{cX}{M_s}} + \tilde{\Lambda}^3 e^{\frac{dAB}{M_s^2}}
\]  

(5.2)

where \( c, d \) will depend on \( a, b \). Expanding the second and third terms to leading order, we find the superpotential (1.1) plus higher-order corrections with \( m = \frac{d\Lambda^3}{M_s^2} \) and \( \mu = \frac{c\Lambda^3}{M_s} \).

Assume that \( h\mu < m^2 \) (which may take some doing). At this order, \( A, B \) have masses of order \( m \). \( X \) is stabilized by a one-loop mass of order \( m_C W = \frac{h^4 m^2}{m^2} < \mu \) (if \( h < 1 \)). Higher order corrections typically render these vacua metastable \([9-11,17,2]\). Furthermore, the F-term for \( X \) is nonvanishing and is the order parameter for SUSY breaking. If \( X \) is a geometric modulus, then this fits within the class of examples discussed in \([12-16]\), in which the F-terms for the moduli provide the dominant contributions to the soft SUSY-breaking terms (we are not addressing here the flavor problems which typically arise in these scenarios – cf. \([27]\) – as well as in gravity mediation). Indeed, much of the literature on gaugino condensation in string theory makes use of this kind of scenario.

Acknowledgements

I would like to thank Michael Dine, Shamit Kachru, John McGreevy, Erich Poppitz, and Howard Schnitzer for discussions and for reading over a draft of this manuscript. I would also like to thank Micha Berkooz, Oliver DeWolfe, Louis Leblond, Martin Schmaltz,

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6 In particular, there will be a term \( \epsilon X^2 \); this will destabilize the SUSY-breaking pseudomoduli space at tree level, apparently causing the system to roll to the supersymmetric point; however, the Coleman-Weinberg potential will lead to a metastable SUSY-breaking vacuum near the origin, so long as \( \epsilon \) is sufficiently small.
Amit Sever, and Eva Silverstein for helpful discussions. Much of this work was completed at the Aspen Center for Physics during the summer 2008 workshop on "Supersymmetry breaking and its mediation in field theory and string theory." I would like to thank the ACP staff and the organizers of the workshop for creating a stimulating and productive working environment. This work was supported by DOE Grant No. DE-FG02-92ER40706, and by a DOE Outstanding Junior Investigator award.
References

[1] K. Intriligator, N. Seiberg and D. Shih, “Supersymmetry Breaking, R-Symmetry Breaking and Metastable Vacua,” JHEP 0707, 017 (2007) [arXiv:hep-th/0703281].

[2] K. Intriligator and N. Seiberg, “Lectures on Supersymmetry Breaking,” Class. Quant. Grav. 24, S741 (2007) [arXiv:hep-ph/0702069].

[3] O. Aharony, S. Kachru and E. Silverstein, “Simple Stringy Dynamical SUSY Breaking,” Phys. Rev. D 76, 126009 (2007) [arXiv:0708.0493 [hep-th]].

[4] M. Gomez-Reino and C. A. Scrucca, “Constraints for the existence of flat and stable non-supersymmetric vacua in supergravity,” JHEP 0609, 008 (2006) [arXiv:hep-th/0606273].

[5] M. Gomez-Reino and C. A. Scrucca, “Metastable supergravity vacua with F and D supersymmetry breaking,” JHEP 0708, 091 (2007) [arXiv:0706.2785 [hep-th]].

[6] S. P. de Alwis, “Effective potentials for light moduli,” Phys. Lett. B 626, 223 (2005) [arXiv:hep-th/0506266].

[7] S. P. de Alwis, “On integrating out heavy fields in SUSY theories,” Phys. Lett. B 628, 183 (2005) [arXiv:hep-th/0506267].

[8] A. Achucarro, S. Hardeman and K. Sousa, “Consistent Decoupling of Heavy Scalars and Moduli in N=1 Supergravity,” [arXiv:0806.4364 [hep-th]].

[9] M. Dine, J. L. Feng and E. Silverstein, “Retrofitting O’Raifeartaigh models with dynamical scales,” Phys. Rev. D 74, 095012 (2006) [arXiv:hep-th/0608159].

[10] M. Dine and J. D. Mason, “Dynamical Supersymmetry Breaking and Low Energy Gauge Mediation,” [arXiv:0712.1353 [hep-ph]].

[11] M. Dine and J. Mason, “Gauge mediation in metastable vacua,” Phys. Rev. D 77, 016005 (2008) [arXiv:hep-ph/0611312].

[12] V. S. Kaplunovsky and J. Louis, “Model independent analysis of soft terms in effective supergravity and in string theory,” Phys. Lett. B 306, 269 (1993) [arXiv:hep-th/9303040].

[13] J. A. Casas, Z. Lalak, C. Munoz and G. G. Ross, “Hierarchical Supersymmetry Breaking and Dynamical Determination of Compactification Parameters by Nonperturbative Effects,” Nucl. Phys. B 347, 243 (1990).

[14] B. de Carlos, J. A. Casas and C. Munoz, “Supersymmetry breaking and determination of the unification gauge coupling constant in string theories,” Nucl. Phys. B 399, 623 (1993) [arXiv:hep-th/9204012].

[15] A. Brignole, L. E. Ibanez and C. Munoz, “Towards a theory of soft terms for the supersymmetric Standard Model,” Nucl. Phys. B 422, 125 (1994) [Erratum-ibid. B 436, 747 (1995)] [arXiv:hep-ph/9308271].

[16] A. Brignole, L. E. Ibanez, C. Munoz and C. Scheich, “Some Issues In Soft Susy Breaking Terms From Dilaton / Moduli Sectors,” Z. Phys. C 74, 157 (1997) [arXiv:hep-ph/9508255].
[17] K. Intriligator, N. Seiberg and D. Shih, “Dynamical SUSY breaking in meta-stable vacua,” JHEP 0604, 021 (2006) [arXiv:hep-th/0602239].

[18] M. Cvetic and T. Weigand, “A string theoretic model of gauge mediated supersymmetry beaking,” arXiv:0807.3953 [hep-th].

[19] Z. Sun, “Continuous degeneracy of non-supersymmetric vacua,” arXiv:0807.4000 [hep-th].

[20] S. Kachru, J. McGreevy and P. Svrcek, “Bounds on masses of bulk fields in string compactifications,” JHEP 0604, 023 (2006) [arXiv:hep-th/0601111].

[21] T. Banks, “Remarks on M theoretic cosmology,” arXiv:hep-th/9906126.

[22] T. Banks, “M-theory and cosmology,” arXiv:hep-th/9911067.

[23] C. Vafa, “Superstrings and topological strings at large N,” J. Math. Phys. 42, 2798 (2001) [arXiv:hep-th/0008142].

[24] A. Lawrence and J. McGreevy, “Local string models of soft supersymmetry breaking,” JHEP 0406, 007 (2004) [arXiv:hep-th/0401034].

[25] A. Lawrence and J. McGreevy, “Remarks on branes, fluxes, and soft SUSY breaking,” arXiv:hep-th/0401233.

[26] A. Lawrence, T. Sander, M. B. Schulz and B. Wecht, “Torsion and Supersymmetry Breaking,” arXiv:0711.4787 [hep-th].

[27] J. Louis and Y. Nir, “Some phenomenological implications of string loop effects,” Nucl. Phys. B 447, 18 (1995) [arXiv:hep-ph/9411429].