Weak Scale Supersymmetry
Without Weak Scale Supergravity

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Abstract

It is generally believed that weak scale supersymmetry implies weak scale supergravity, in the sense that the masses of the gravitino and gravitationally coupled moduli have masses below 100 TeV. This paper presents a realistic framework for supersymmetry breaking in the hidden sector in which the masses of the gravitino and gravitational moduli can be much larger. This cleanly eliminates the cosmological problems of hidden sector models. Supersymmetry breaking is communicated to the visible sector by anomaly-mediated supersymmetry breaking. The framework is compatible with perturbative gauge coupling unification, and can be realized either in models of ‘warped’ extra dimensions, or in strongly-coupled four-dimensional conformal field theories.
Supersymmetry (SUSY) is arguably the most attractive framework for explaining the origin of electroweak symmetry breaking. SUSY automatically stabilizes the weak scale against quantum corrections (the ‘hierarchy problem’) and is naturally compatible with the absence of large corrections in precision electroweak data. If SUSY exists in nature it must be broken, and understanding the possible origin of SUSY breaking and its implications for future experiments is one of the central problems in particle physics.

The only truly model-independent prediction of SUSY is the existence of superpartners of the observed particles, plus spin 0 Higgs particles and their superpartners. If SUSY solves the hierarchy problem, then the masses of squarks, sleptons, gauginos, Higgs fields, and Higgsinos must be below of order 1 TeV.

In this paper, we will focus on the gravitino, the spin \( \frac{3}{2} \) superpartner of the graviton. It is commonly thought that the gravitino mass cannot be more than of order 100 TeV in models where SUSY solves the hierarchy problem, as we now explain. SUSY breaking gives rise to a gravitino mass of order

\[
m_{3/2} \sim \frac{F}{M_P},
\]

where \( F \) is the SUSY breaking order parameter and \( M_P \sim 10^{18} \text{ GeV} \) is the Planck scale. The size of \( F \) is model dependent, and depends on the strength of the ‘messenger’ interactions that communicate SUSY breaking to the visible sector. Given the fact that SUSY breaking masses in the visible sector are between 100 GeV and 1 TeV, the gravitino mass can be large if the messenger interactions are weak. Gravity necessarily couples the visible and hidden sectors with universal strength, and is therefore the weakest possible messenger of SUSY breaking. In general models of gravity mediated SUSY breaking, gravitational contact terms give rise to superpartner masses of order \( m_{3/2} \), so in these theories \( m_{3/2} \lesssim 1 \text{ TeV} \).

It is also possible to suppress the contact interactions between the visible and hidden sectors [1, 2]. In such models, the supergravity contribution to supersymmetry breaking is related to the conformal anomaly [1, 3]. The ‘anomaly mediated’ superpartner masses are of order \((g_{SM}^2/16\pi^2)m_{3/2}\), which implies \( m_{3/2} \lesssim 100 \text{ TeV} \).

String theory and higher-dimensional supergravity also generically predict the existence of numerous moduli fields with Planck suppressed couplings to visible matter. In the presence of SUSY breaking, these generally get masses of order \( F/M_P \sim m_{3/2} \), so we expect \( m_{\text{moduli}} \lesssim 100 \text{ TeV} \).

In this paper, we show that the gravitino and moduli masses can naturally be much larger. This is interesting because it gives a general solution to the cosmological problems associated with the gravitino and moduli. Gravitinos and moduli
are readily produced in inflationary reheating, and live long because of their Planck-suppressed couplings to matter. If the gravitino decays during or after nucleosynthesis, it can upset the predictions of nucleosynthesis. For $m_{3/2} \gtrsim 60 \text{ TeV}$ the gravitino decays sufficiently rapidly to avoid this problem \[4\]. For smaller gravitino mass, we can require the reheat temperature after inflation to be low enough to suppress the gravitino abundance. Refs. [5] obtain $T_{\text{reheat}} \lesssim 10^8 \text{ GeV}$ for $m_{3/2} \sim 1 \text{ TeV}$, but recent work suggests that the gravitino production is more efficient, requiring much lower reheat temperatures [6]. Similar bounds apply to moduli, but the bounds are more model-dependent.

For spin 0 moduli fields, there is also the ‘Polonyi problem’ \[7\]. Briefly stated, the values of the moduli fields in the early universe differ from their present vacuum values. This stores energy, and when this energy is released it generally reheats the universe to a temperature too low for successful nucleosynthesis. For $m_{\text{modulus}} \gtrsim 100 \text{ TeV}$, the reheat temperature is $\gtrsim 1 \text{ MeV}$, which is just enough for nucleosynthesis.

To summarize the cosmological bounds, it is fair to say that anomaly mediation narrowly satisfies the bounds. However, it is important that there is a class of models in which the gravitino and moduli masses are larger and the bounds are satisfied by a wide margin.

We now show how this can be realized in the supersymmetric Randall–Sundrum model \[8\]. This is a 5D effective field theory where the 5th dimension is compactified on an interval of length $\pi r$, realized as a $S^1/Z_2$ orbifold. The metric can be written

$$ds^2 = e^{-2k|\vartheta|} \eta_{\mu\nu} dx^\mu dx^\nu + r^2 d\vartheta^2, \quad -\pi < \vartheta \leq +\pi.$$  \hspace{1cm} (2)

The slope discontinuities in the metric at $\vartheta = 0$, $\pi$ are due to the presence of 4D branes that are fixed at the boundary by orbifold boundary conditions. The presence of the ‘warp factor’ $e^{-2k|\vartheta|}$ in the metric means that all physical scales on the ‘IR brane’ at $\vartheta = \pi$ are redshifted compared to the scales on the ‘UV brane’ at $\vartheta = 0$. This model is supersymmetric with the addition of appropriate additional fields and interactions \[9\].

At energies below the mass of the lightest gravitational Kaluza–Klein (KK) mode the theory can be described by a 4D effective lagrangian consisting of 4D SUGRA coupled to the radion \[10, 11\]:

$$\mathcal{L}_{4, \text{eff}} = -\frac{3M_5^3}{k} \int d^4 \theta \left( \phi^+ \phi - \omega^+ \omega \right) + \int d^4 \theta \left( \phi^+ \phi K_{\text{UV}} + \omega^+ \omega K_{\text{IR}} \right)$$  

$$+ \left[ \int d^2 \theta \left( \phi^3 W_{\text{UV}} + \omega^3 W_{\text{IR}} \right) + \text{h.c.} \right].$$  \hspace{1cm} (3)
Here $\phi$ is the conformal compensator, and the ‘warp factor’ superfield

$$\omega = \phi e^{-kT}, \quad T = \pi r + \cdots$$

parameterizes the radion. The first term contains the SUGRA and radion kinetic terms, and the remaining terms come from Kähler potentials and superpotentials localized on the UV and IR branes.

We will be interested in the scenario where the visible sector fields are localized on the IR brane and SUSY is broken on the UV brane. Provided that there are no additional light fields in the bulk, this automatically suppresses flavor-violating contact terms between the visible and hidden sectors, ‘sequestering’ the hidden sector [1]. This naturally explains the absence of flavor-changing neutral currents from superpartners. We will be interested in the anomaly-mediated SUSY breaking (AMSB) contribution in the visible sector. The regulator for loops of visible sector fields must be localized on the IR brane. The conformal compensator for these loops is therefore $\omega$, and anomaly-mediated masses are of order $(g_{\text{SM}}^2/16\pi^2)M_{\text{AMSB}}$, with $M_{\text{AMSB}} = \langle F_\omega/\omega \rangle$. The size of $M_{\text{AMSB}}$ depends on the mechanism for SUSY breaking and radius stabilization.

Radius stabilization for 5D supergravity theories in a consistent effective field theory framework was first achieved in Refs. [12, 10]. However, the stabilization mechanism of Ref. [10] gives $\langle F_\omega/\omega \rangle \sim \langle F_\phi \rangle$. We consider a stabilization sector that gives rise to the following additional terms in the 4D effective lagrangian:

$$\triangle L_{\text{4, eff}} = \int d^2 \theta \left( c_{\text{UV}} \phi^3 + c_{\text{IR}} \omega^3 + \epsilon \phi^{3-n} \omega^n \right) + \text{h.c.}$$

$$- F_{\text{UV}}^2 \left[1 + \text{Goldstino terms} \right].$$

The first two terms can arise from constant superpotentials localized on the UV and IR branes, respectively; the 5D origin of the third term will be described below; the last term represents the effect of SUSY breaking on the UV brane.

In the limit $\epsilon \to 0$, the vacuum is at $\langle \omega \rangle = 0$, corresponding to infinite separation between the UV and the IR branes. In this limit the fundamental scale on the IR brane is $M_{\text{IR}} = M_5 \omega$, and the observations of an IR brane observer are defined relative to this scale. For example, the mass of the lowest KK mode is $m_{\text{KK}} \sim k\omega$, which is proportional to $M_{\text{IR}}$, so an IR observer sees a finite mass gap in the KK spectrum. However, the 4D Planck scale is $M_\text{P}^2 \sim M_{\text{IR}}^2/\omega^2 \to \infty$, so 4D gravity is decoupled from physics on the IR brane. In this limit it is easy to read off the SUSY breaking from Eq. (5) because there is no mixing between the $\omega$ and $\phi$ fields, and hence no supergravity corrections to the $\omega$ potential. The scale of AMSB is given by $F_\omega/\omega \sim \omega$,
which is also proportional to $M_{\text{IR}}$. In this limit an observer on the IR brane sees SUSY broken by anomaly mediation, even though gravity has completely decoupled!\(^1\) This magic is due to conformal invariance. In the limit $\epsilon \to 0$ the terms that depend on $\omega$ have an exact (nonlinearly realized) conformal symmetry. This ensures that they are independent of the conformal compensator $\phi$.

To get a model with 4D gravity, we want $\langle \omega \rangle \neq 0$. For $\epsilon \ll c_{\text{UV,IR}}$ and $n < 3$, the $\epsilon$ term gives a small shift to the vacuum:

$$|\langle \omega \rangle|^{1-n} = \frac{n(3-n)}{6} \left| \frac{\epsilon c_{\text{UV}}}{c_{\text{IR}}^2} \right| \ll 1$$

with

$$\left| \frac{\langle F_\omega \rangle}{\langle \omega \rangle} \right| = \left| \frac{c_{\text{IR}}}{M_P^2} \langle \omega \rangle \right|, \quad |\langle F_\phi \rangle| = \left| \frac{c_{\text{UV}}}{M_P^2} = \frac{F_{\text{UV}}}{\sqrt{3}M_P} \right|,$$

where $M_P^2 = M_3^2/k$. The radion mass is of order $\langle F_\omega/\omega \rangle$, while the mass of the gravitino is of order $\langle F_\phi \rangle$. Bulk moduli fields will also have a mass of order $\langle F_\phi \rangle$ provided that the wavefunction of the lightest KK mode has sizable overlap with the UV brane. Note that the order parameter for SUSY breaking on the IR brane is parametrically suppressed (by $\langle \omega \rangle \ll 1$) compared to $\langle F_\phi \rangle$.

We now describe the bulk interactions that give rise to the $\epsilon$ term in Eq. (5). We add a $SU(2)$ gauge multiplet in the bulk, with 6 fundamentals with mass $m$ localized on the UV brane. Below the scale $m_{\text{KK}}$ this becomes a 4D $SU(2)$ gauge theory with 6 fundamentals. If $m \lesssim m_{\text{KK}}$, this theory generates a dynamical superpotential of the form Eq. (5), with

$$\epsilon \sim \left( \frac{m}{g_5^2} \right)^{3/2}, \quad n = \frac{4\pi^2}{kg_5^2},$$

where $g_5$ is the 5D gauge coupling. The small $m$ condition gives the constraint

$$m_{3/2} \gtrsim \frac{M_P^2 M_{\text{AMSB}}^2}{k^3 \langle \omega \rangle^{n-1/2}},$$

where we use $g_5^2 \sim 1/k^2$ for $n \sim 1$. There are also contact terms between the $SU(2)$ fundamentals and the SUSY breaking sector. These give rise to corrections to the radion potential of order $\Delta V \sim m_{3/2}^2 (\epsilon \omega^n)^{3/2}$, which can be neglected provided

$$m_{3/2}^2 \lesssim M_{\text{AMSB}} m_{\text{KK}}.$$\(^2\)

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\(^1\)I thank R. Sundrum for a discussion of this point.
The inequalities Eqs. (9) and (10) justify the use of the 4D effective field theory above for stabilizing the radion.

We now investigate the conditions under which \( \langle F_\omega / \omega \rangle \) dominates SUSY breaking for fields on the IR brane. First we consider SUSY breaking from SUGRA loops. In a 4D effective theory, the loop is dominated by a quadratically divergent contribution from 4D momenta above \( m_{3/2} \). (We are restricting attention to the case \( m_{3/2} \lesssim m_{\text{KK}} \) where we can treat the gravitino in the 4D effective theory.) This gives

\[
\Delta m^2_{\text{scalar}} \sim \frac{m_{3/2}^2 \Lambda_{4D}^2}{16 \pi^2 M_P^4},
\]

where \( \Lambda_{4D} \) is the UV cutoff in the 4D theory. In the 5D theory the integral is cut off because in position space the SUGRA propagators must extend from the IR brane to the UV brane in order to communicate SUSY breaking. Because the gravitino loop cannot shrink to zero size, so there is no UV divergence [1]. We can understand the size of this finite loop contribution as follows. For 4D loop momenta \( p_4 \lesssim m_{\text{KK}} \), the graviton loop behaves as in the 4D theory. For \( p_4 \gtrsim m_{\text{KK}} \) the brane-to-brane propagator falls off as \( e^{-|p_4|/m_{\text{KK}}} \), so the result is given by Eq. (11) with \( \Lambda_{4D} \sim m_{\text{KK}} \).

The AMSB contribution is larger than the gravity loop contribution provided that

\[
\frac{m_{3/2}}{M_{\text{AMSB}}} \lesssim \frac{M_P}{4 \pi m_{\text{KK}}}.
\]

SUSY breaking can also be communicated to the visible sector through loops of \( SU(2) \) gauge fields from the stabilization sector. The \( SU(2) \) gauge fields couple to observable fields via flavor-violating interactions on the IR brane of the form

\[
\Delta \mathcal{L}_{\text{IR}} \sim \int d^4 \theta \omega^4 \frac{1}{M_p^2} Q^1 Q W^\alpha \sigma^\mu \partial_\mu \bar{W}^\alpha,
\]

where \( Q \) is a visible sector field and \( W_\alpha \) is the \( SU(2) \) field strength.\(^3\) The \( SU(2) \) gaugino gets a mass from the coupling to the radion of order \( \langle F_T / T \rangle \sim \langle F_\phi / \ln \omega \rangle \) [14], which gives rise to flavor-violating visible scalar masses. The condition that these are smaller than experimental bounds gives

\[
\frac{m_{3/2}}{M_{\text{AMSB}}} \lesssim 10^{-2} \left( \frac{M_p}{k} \right)^{5/3} \left( \frac{\ln \omega}{\omega} \right)^{5/2}.
\]

\(^2\)The case where this inequality is saturated may be interesting. Ref. [13] did the calculation for a flat extra dimension and obtained a negative mass-squared. The calculation has not been done for the warped case.

\(^3\)In Eq. (13) we define all fields to have vanishing conformal weight.
Provided that Eqs. (9), (10), (12), and (14) are satisfied, SUSY breaking for fields localized on the IR brane is dominated by anomaly-mediated SUSY breaking from $\langle F_\omega/\omega \rangle$.

We now discuss the phenomenology of this model. SUSY breaking is communicated to the visible sector by anomaly mediation. This naturally gives a flavor-blind squark masses and therefore explains the absence of flavor-changing neutral currents from squark mixing. If the visible sector is the minimal supersymmetric standard model, the slepton mass-squared terms are negative [1]. There are a number of proposals in the literature for non-minimal models that give a realistic spectrum while preserving the attractive features of anomaly mediation [15].

In order to preserve the success of perturbative gauge coupling unification, we require that $M_5 \omega \gtrsim M_{GUT} \sim 10^{16}$ GeV. The largest allowed value for the gravitino mass in this scenario is $m_{3/2} \sim 10^9$ GeV, with $m_{KK} \lesssim 10^{13}$ GeV and $\langle \omega \rangle \sim 10^{-2}$. For $k \sim M_5$ the largest value is $m_{3/2} \sim 10^6$ GeV. If one is willing to give up perturbative gauge coupling unification, higher values of $m_{3/2}$ can be achieved.

The duality between 5D anti De Sitter space and 4D conformal field theory (CFT) [16, 17] suggests that there is another realization of this scenario where the role of the 5D warped bulk is played by a strongly coupled 4D CFT. Ref. [2] showed that 4D CFT’s with no known 5D ‘dual’ description can lead to a sequestered hidden sector and anomaly mediated SUSY breaking. Along these lines, we now show that the qualitative features of the 5D models described above can be realized in a 4D theory based on a strongly-coupled CFT. This gives additional insight into the mechanism described above.

Consider a 4D $SU(2)$ SUSY gauge theory with 8 fundamentals $P$. This theory is asymptotically free in the UV, and approaches a strongly coupled conformal fixed point in the IR [18]. We add to this theory a superpotential

$$W = \lambda P^4 + \kappa P^2. \tag{15}$$

At the fixed point, the dimension of $P$ is $\frac{3}{4}$, so $\lambda$ is dimensionless and $\kappa$ has dimension $\frac{3}{2}$. We can parameterize the moduli space by

$$PP = \begin{pmatrix} 2 & 6 \\ 6 & 0 \end{pmatrix} \begin{pmatrix} 0 & X \\ -X & 0 \end{pmatrix} \begin{pmatrix} Y \\ -Y^T \end{pmatrix}, \tag{16}$$

and expand about $X \neq 0$, $Y = 0$. The (nonlinearly realized) conformal symmetry
implies that the effective field theory below the scale $X$ has Kähler terms [2]

\[
\Delta L_{\text{eff}} = \int d^4 \theta \phi^\dagger \phi (X^\dagger X)^{2/3} \left[ 1 + \mathcal{O}(|Y|^2/|X|^2) \right] \\
= \int d^4 \theta \hat{X}^\dagger \hat{X} \left[ 1 + \mathcal{O}(|\hat{Y}|^2/|\hat{X}|^2) \right],
\]

(17)

(18)

where $\hat{X} = \phi X^{2/3}$, $\hat{Y} = \phi Y/X^{1/3}$. The $\phi$ dependence can be completely scaled away because of conformal invariance. The effective superpotential terms are

\[
\Delta L_{\text{eff}} = \int d^2 \theta \left[ \lambda \hat{X}^3 + \epsilon \phi^{3-n} \hat{X}^n + (Y \text{ dependent}) \right] + \text{h.c.},
\]

(19)

with $n = \frac{3}{2}$. This is precisely the same effective lagrangian we obtained in the 5D case with the identifications $\hat{X} \leftrightarrow M_P \omega$, $\lambda \leftrightarrow c_{\text{IR}}/M_P^3$, $\kappa \leftrightarrow \epsilon/M_P^n$.

The composite fields $\hat{Y}$ correspond to the fields localized on the IR brane in the 5D model. Loops of $\hat{Y}$ fields are cut off at a scale proportional to $\hat{X}$, so the anomaly-mediated contribution to their masses is determined by $\langle F_X / \hat{X} \rangle \ll \langle F_\phi \rangle$, as in the 5D model. The suppression of anomaly-mediated contribution proportional to $F_\phi$ in both frameworks is closely related to conformal invariance.

A complete model requires that the standard model gauge bosons are composite. This is compatible with perturbative unification if the compositeness scale is above the GUT scale. Sequestering also requires that there are no unbroken global (e.g. flavor) symmetries in the visible sector [2]. Ref. [19] gives an interesting mechanism for generating the observed flavor structure in this type of scenario. We conclude that the ingredients for a fully realistic model can be realized in 4D CFT’s. Our ability to construct explicit realistic 4D models is limited mainly by our poor understanding of strongly coupled superconformal field theories.

We have restricted attention to the regime $m_{3/2} \lesssim m_{\text{KK}}$, where the gravitino can be treated in the 4D effective field theory. There appears to be no fundamental reason that we cannot have $m_{3/2} \gg m_{\text{KK}}$, and this is presently under investigation.

In conclusion, we have shown that the masses of the gravitino and gravitational moduli can be much larger than the weak scale, without upsetting the major successes of supersymmetry: a solution to the hierarchy problem and the successful predictions of perturbative gauge coupling unification. The only assumption about the moduli is that SUSY breaking in the moduli has gravitational strength. This may be of particular interest in superstring/M theory, where quasi-realistic compactifications have many moduli that give rise to cosmological difficulties.

\[\text{Composite gauge bosons are known to emerge in simple SUSY gauge theories [18].}\]
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References


