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Gaugino mediated supersymmetry breaking

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Abstract: We consider supersymmetric theories where the standard-model quark and lepton fields are localized on a '3-brane' in extra dimensions, while the gauge and Higgs fields propagate in the bulk. If supersymmetry is broken on another 3-brane, supersymmetry breaking is communicated to gauge and Higgs fields by direct higher-dimension interactions, and to quark and lepton fields via standard-model loops. We show that this gives rise to a realistic and predictive model for supersymmetry breaking. The size of the extra dimensions is required to be of order 10–100 times larger than fundamental scale (e.g. the string scale). The spectrum is similar to (but distinguishable from) the predictions of 'no-scale' models. Flavor-changing neutral currents are naturally suppressed. The μ term can be generated by the Giudice-Masiero mechanism. The supersymmetric CP problem is naturally solved if CP violation occurs only on the observable sector 3-brane. These are the simplest models in the literature that solve all supersymmetric naturalness problems.

Keywords: Field Theories in Higher Dimensions, Supersymmetry Breaking, Supersymmetric Standard Model.

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Contents

1. Introduction

Supersymmetry (SUSY) provides an attractive framework for solving the hierarchy problem, but it introduces naturalness puzzles of its own. Perhaps the most serious is the 'SUSY flavor problem:' why do the squark masses conserve flavor? A natural solution is given by models of gauge-mediated SUSY breaking [1] or by 'anomalous U(1)' models [2]. In ref. [3], Randall and Sundrum suggested another solution in theories where the visible sector fields are localized on a '3-brane' in extra dimensions and the hidden sector fields are localized on a spatially separated '3-brane'. (Models of this type were introduced in the context of string theory by Hořava and Witten [4].) Ref. [3] pointed out that in such theories contact terms between the visible and hidden fields are suppressed if the separation r between the visible and hidden branes is sufficiently large. The reason is simply that contact terms arising from integrating out states with mass M are suppressed by a Yukawa factor e^{-Mr} if $M \gtrsim r$. Because the suppression is exponential, the separation need only be an order of magnitude larger than the fundamental scale (e.g. the string scale) to strongly suppress contact interactions.

If contact interactions between the hidden and visible sector fields can be neglected, other effects become important for communicating SUSY breaking. One possibility is the recently-discovered mechanism of anomaly-mediation [3, 5], a model-independent supergravity effect that is always present. (For a careful discussion of anomaly mediation in a specific higher-dimensional model, see ref. [6].) Unfortunately, if anomaly-mediation dominates, and if the visible sector is the minimal supersymmetric standard model (MSSM), then slepton mass-squared terms are negative. This problem can be avoided in extensions of the MSSM [7]. In this paper, we will explore the alternate possibility that standard-model gauge and Higgs fields propagate in the bulk and communicate SUSY breaking between the hidden sector and visible-sector matter fields. Models with the all the MSSM superfields except the gauge and Higgs fields localized on a 3-brane were also considered in ref. [8]. In those models supersymmetry was directly broken by the compactification boundary conditions, requiring a rather large extra dimension (radius of order TeV^{-1}) in order to explain the gauge hierarchy. Models similar to the one considered here , i.e. with a hidden supersymmetry breaking sector sequestered on a different 3-brane and standard-model gauge and Higgs fields in the bulk, were considered in ref. [9]. These models contained an additional U(1) gauge multiplet; the present paper shows that this is not required to obtain a realistic theory of SUSY breaking.

In the higher-dimensional theory, the standard-model gauge and Higgs fields can interact only through non-renormalizable interactions. We therefore treat the higher-dimensional theory as an effective theory with a cutoff M, which may be viewed as the fundamental scale of the theory. Below the compactification scale $\mu_c \sim 1/r$, the theory matches onto a 4-dimensional effective theory. In this theory, the couplings of the gauge and Higgs fields are suppressed by $1/(Mr)^{D-4}$, where D is the number of 'large' spacetime dimensions. Therefore, the size of the extra dimensions cannot be too large in units of the fundamental scale. However, because the suppression of contact terms is exponential, there is a range of radii with sufficient suppression of contact terms to avoid flavor-changing neutral currents without exceeding the strong-coupling bounds on the couplings in the higher-dimension theory [9].

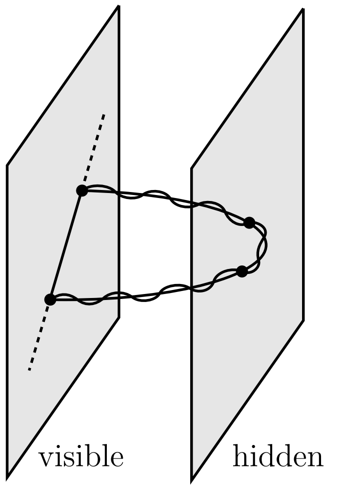

Figure 1: The leading diagram that contributes to SUSY-breaking scalar masses in the models considered in this paper. The bulk line is a gaugino propagator with two mass insertions on the hidden brane.

In this scenario, SUSY breaking masses for gauginos and Higgs fields are generated by higher-dimension contact terms between the bulk fields and the hidden sector fields, assumed to arise from a more fundamental theory such as string theory. In particular, the μ term can be generated by the Giudice-Masiero mechanism [10]. Other direct contact interactions between the hidden and visible sectors are suppressed because of their spatial separation. The leading contribution to SUSY breaking for visible sector fields arises from loops of bulk gauge and Higgs fields, as illustrated in figure 1. These diagrams are ultraviolet convergent (and hence calculable) because the spatial separation of the hidden and visible branes acts as a physical point-splitting regulator. In effective field theory language, the contribution from loop momenta above the compactification scale is a (finite) matching contribution, while the contribution from loop momenta below the compactification scale can be obtained from the 4-dimensional effective theory. The higher-dimensional theory therefore gives initial conditions for the 4-dimensional renormalization group at the compactification scale μ_c: nonzero gaugino masses and Higgs mass parameters, and loop-suppressed soft SUSY breaking parameters for the squarks and sleptons. This is similar to the boundary conditions of 'no-scale' supergravity models [11], but in the present case the boundary conditions are justified by the geometry of the higher-dimensional theory. Since the SUSY breaking masses for all chiral matter fields other than the third generation squarks are dominated by the gaugino loop, we call this scenario 'gaugino mediated SUSY breaking' ($\tilde{g}$MSB).

The renormalization group has a strong effect on the SUSY breaking parameters, and the soft masses at the weak scale are all of the same order. In fact, the Bino can be the lightest superpartner (LSP) in this scenario. The spectrum is similar to that of 'no-scale' supergravity models [11], with the important difference that the present scenario allows a Fayet-Iliopoulos term for hypercharge that can have an important effect on the slepton spectrum. We obtain realistic spectra without excessive fine-tuning for neutralino and slepton masses below approximately 200 GeV, suggesting that these superpartners are relatively light in this scenario.

This paper is organized as follows. In section 2, we discuss the higher-dimensional theory. We show that the size of the extra dimensions can be large enough to suppress FCNC's while still having gauge and Higgs couplings of order 1 at low energies. We also show how the SUSY CP problem can be naturally solved in this class of models. In section 3, we discuss the phenomenology of this class of models. section 4 contains our conclusions.

2. Bulk gauge and Higgs fields

In this section, we discuss some general features of higher-dimensional theories with gauge and Higgs fields in the bulk and other fields localized on '3-branes.' We use the term '3-branes' to mean either dynamical surfaces (e.g. topological defects or string-theory D-branes) or non-dynamical features of the higher-dimensional spacetime (e.g. orbifold fixed points). All of these ingredients occur in string theory, but we will not concern ourselves with the derivation of the model from a more fundamental theory. We simply write an effective field theory valid below some scale M, which may be the string scale, the compactification scale associated with additional small dimensions, or some other new physics.

We therefore consider an effective theory with D spacetime dimensions, with $3+1$ non-compact spacetime dimensions and $D-4$ compact spatial dimensions with linear size of order r. The D-dimensional effective lagrangean takes the form

$$\mathcal{L}_D = \mathcal{L}_{\text{bulk}}(\Phi(x,y)) + \sum_j \delta^{D-4}(y-y_j)\mathcal{L}_j(\Phi(x,y_j),\phi_j(x))\,, \tag{2.1}$$

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where j runs over the various branes, x are coordinates for the 4 non-compact spacetime dimensions, y are coordinates for the $D-4$ compact spatial dimensions, Φ is a bulk field, and ϕ_j is a field localized on the j^{th} brane. This effective theory can be treated using the usual techniques of effective field theory, and parameterizes the most general interactions of the assumed degrees of freedom below the scale M.[1]

We assume that the $D-4$ extra spatial dimensions are compactified on a distance of order $r \gg 1/M$. We also assume that the distance between different branes is also of order r. This ensures that contact interactions between fields on different branes arising from states above the cutoff are suppressed by the Yukawa factor e^{-Mr}.

We assume that the standard-model gauge and Higgs fields propagate in the bulk. Bulk gauge fields have a gauge coupling with mass dimension $4-D$, which is an irrelevant interaction for all $D>4$. When we match onto the 4-dimensional theory at the compactification scale, the effective 4-dimensional gauge coupling is[2]

$$g_4^2 = \frac{g_D^2}{V_{D-4}}\,, \tag{2.2}$$

where g_D is the gauge coupling in the D-dimensional theory and $V_{D-4} \sim r^{D-4}$ is the volume of the compact dimensions. If $g_D \sim 1/M^{D-4}$, we have $g_4 \sim 1/(Mr)^{D-4} \ll 1$, which is unacceptable. In order to have $g_4 \sim 1$ (as observed), we require the gauge coupling to be larger than unity in units of M. However, it presumably does not make sense to take g_D larger than its strong-coupling value, defined to be the value where loop corrections are order 1 at the scale M. This follows from 'naïve dimensional analysis' (NDA) [14, 15], which is known to work extremely well in supersymmetric theories [16]. If we assume that the loop corrections are suppressed by ϵ at the scale M, the lagrangean is [9]

$$\mathcal{L}_D \sim \frac{M^D}{\epsilon \ell_D}\mathcal{L}_{\text{bulk}}\left(\frac{\hat{\Phi}}{M}, \frac{\partial}{M}\right) + \sum_j \delta^{D-4}(y-y_j)\frac{M^4}{\epsilon \ell_4}\mathcal{L}_j\left(\frac{\hat{\Phi}}{M}, \frac{\hat{\phi}_j}{M}, \frac{\partial}{M}\right), \tag{2.3}$$

where $\ell_D = 2^D \pi^{D/2}\Gamma(D/2)$ is the geometrical loop factor for D dimensions, and all couplings in $\mathcal{L}_{\text{bulk}}$ and $\mathcal{L}_j$ are order 1. Note that the fields $\hat{\Phi}$ and $\hat{\phi}$ in eq. (2.3) do not have canonical kinetic terms. The idea behind eq. (2.3) is that the factors multiplying $\mathcal{L}_{\text{bulk}}$ and $\mathcal{L}_j$ act as loop-counting parameters (like $\hbar$ in the semiclassical expansion) that cancel the loop factors and ensure that loop corrections are suppressed by ϵ. Strong coupling corresponds to $\epsilon \sim 1$.

We can use eq. (2.3) to read off the value of the D-dimensional gauge coupling

$$g_D^2 \sim \frac{\epsilon \ell_D}{M^{D-4}}\,. \tag{2.4}$$

[1]For an explicit supersymmetric example and calculations, see ref. [12].
[2]We neglect the effects of gravitational curvature.

D	Torus: $ML_{\max}$	$e^{-ML_{\max}/2}$	Sphere: $Mr_{\max}$	$e^{-Mr_{\max}}$
5	740	3×10^{-162}	118	4×10^{-52}
6	63	2×10^{-14}	18	2×10^{-8}
7	29	6×10^{-7}	11	3×10^{-5}
8	20	5×10^{-5}	8.7	2×10^{-4}
9	16	3×10^{-4}	8.0	3×10^{-4}
10	14	9×10^{-4}	7.8	4×10^{-4}
11	13	1×10^{-3}	7.8	4×10^{-4}

Table 1: Estimates of the maximum size and exponential suppression factor for propagation between two branes of maximal separation. $L_{\max}$ is the maximum length of a cycle of a symmetric torus, and $r_{\max}$ is the maximum radius of the sphere.

We can obtain the maximum value for the size of the extra dimension consistent with the fact that $g_4 \sim 1$ by setting $\epsilon \sim 1$ and using eq. (2.2). The results are shown in table 1. We see that the exponential suppression factor due to the large size of the extra dimensions can be substantial even for many extra dimensions [9]. Similar conclusions hold for the Higgs interactions.

To see how much suppression is required, note that the dangerous contact terms have the form (using eq. (2.3))

$$\Delta\mathcal{L}_{\text{brane}} \sim \frac{e^{-Mr}}{M^2} \int d^4\theta \, (\hat{\phi}^\dagger_{\text{hid}} \hat{\phi}_{\text{hid}})(\phi^\dagger_{\text{obs}} \phi_{\text{obs}}) \,, \tag{2.5}$$

where the observable fields (but not the hidden fields) have been canonically normalized. This must be compared with the operators that give rise to the gaugino and Higgs SUSY breaking parameters. From eq. (2.3) we obtain

$$\begin{aligned} \Delta\mathcal{L}_{\text{brane}} \sim \frac{\ell_D}{\ell_4} \left(\int d^2\theta \, \frac{1}{M^{D-3}} \hat{\phi}_{\text{hid}} W^\alpha W_\alpha + \text{h.c.} \right) + \\ + \frac{\ell_D}{\ell_4} \int d^4\theta \left\{ \frac{1}{M^{D-3}} \left(\hat{\phi}^\dagger_{\text{hid}} H_u H_d + \text{h.c.} \right) + \right. \\ \left. + \frac{1}{M^{D-4}} \hat{\phi}^\dagger_{\text{hid}} \hat{\phi}_{\text{hid}} \left[H^\dagger_u H_u + H^\dagger_d H_d + (H_u H_d + \text{h.c.}) \right] \right\}, \end{aligned} \tag{2.6}$$

where W_α is the gauge field strength and $H_{u,d}$ are the Higgs fields, normalized to have canonical kinetic terms in D dimensions. (More precisely, these are $\mathcal{N} = 1$ superfields obtained by projecting the bulk supermultiplets onto the branes. For a specific example, see ref. [12].) Matching to the D-dimensional theory, we find

$$m_{1/2}, \mu \sim \frac{\hat{F}_{\text{hid}}}{M} \frac{\ell_D/\ell_4}{M^{D-4} V_{D-4}} \,, \qquad B\mu, m^2_{H_u}, m^2_{H_d} \sim \frac{\hat{F}^2_{\text{hid}}}{M^2} \frac{\ell_D/\ell_4}{M^{D-4} V_{D-4}} \,. \tag{2.7}$$

Note that the $B\mu$ term and the Higgs mass-squared terms are enhanced by a volume factor.[3] For example,

$$\frac{B\mu}{m_{1/2}^2} \sim \frac{\ell_4}{\ell_D} M^{D-4} V_{D-4} \sim \epsilon \ell_4 \,, \tag{2.8}$$

where we have imposed $g_4 \sim 1$ to obtain the last estimate. We see that if the theory is strongly coupled at the fundamental scale, we require a fine tuning of order $1/\ell_4 \sim 1\%$ to obtain all SUSY breaking parameters of the same size [9]. However, for a small number of extra dimensions, the fundamental theory need not be strongly coupled at the fundamental scale, and we can naturally obtain all SUSY breaking terms close to the same size. For example for $D = 5$ compactified on a circle with circumference $L \sim 20/M$, the exponential suppression is $e^{-10} \sim 5 \times 10^{-5}$ and $B\mu/m_{1/2}^2 \sim 4$. As the number of extra dimensions increases, the strong coupling estimate is approached rapidly. See table 2.

D	Torus: $B\mu/m_{1/2}^2$	Sphere: $B\mu/m_{1/2}^2$
5	3.4	3.4
6	10	16
7	28	43
8	69	85
9	160	130

Table 2: Estimates of $B\mu/m_{1/2}^2$ for the symmetric torus and the sphere. The size of the extra dimension is chosen so that the exponential suppression factor is of order $e^{-8} \simeq 3 \times 10^{-4}$ (approximately the maximum suppression for large D). This means that the torus has cycle length $L = 16/M$, and the sphere has radius $r = 8/M$.

The contribution to visible sector scalar masses from contact terms is

$$\Delta m_{\text{vis}}^2 \sim e^{-Mr} \left(\frac{\hat{F}_{\text{hid}}}{M} \right)^2 . \tag{2.9}$$

The values eq. (2.7) are the values renormalized at the compactification scale; we will later see that we require $\hat{F}_{\text{vis}}/M \sim 200\,\text{GeV}$. Using the experimental constraints[4]

$$\frac{m_{\tilde{d}\tilde{s}}^2}{m_{\tilde{s}}^2} \lesssim (6 \times 10^{-3}) \left(\frac{m_{\tilde{s}}}{1\,\text{TeV}} \right), \qquad \text{Im} \left(\frac{m_{\tilde{d}\tilde{s}}^2}{m_{\tilde{s}}^2} \right) \lesssim (4 \times 10^{-4}) \left(\frac{m_{\tilde{s}}}{1\,\text{TeV}} \right), \tag{2.10}$$

so $e^{-Mr} \sim 10^{-3}$ to 10^{-4} is plausibly sufficient to suppress FCNC's.

We now discuss the loop effects that communicate SUSY breaking to the visible sector fields, such as those illustrated in figure 1. These are ultraviolet convergent because the separation of the hidden and visible branes acts as a physical point-splitting regulator for these diagrams. Another way to see this is that there is no *local* counterterm in the D-dimensional theory that can cancel a possible overall divergence.[5] Given a specific D-dimensional theory, this diagram is therefore calculable. From the

[3] This point was missed in an earlier version of this paper. It was pointed out in ref. [13], which appeared while this paper was being completed. See also ref. [9].

[4] For a complete discussion, see e.g. ref. [17].

[5] Multiloop diagrams may have subdivergences, but these can always be cancelled by counterterms localized on one of the branes.

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point of view of 4-dimensional effective field theory, the extra dimensions act as a cutoff of order $\mu_c \sim 1/r$. The effects of this cutoff can be absorbed into a counterterm for the visible sector scalar masses and A terms of order

$$\Delta m^2_{\rm vis} \sim \frac{g_4^2}{16\pi^2} m^2_{1/2}\,, \qquad \Delta A_{\rm vis} \sim \frac{g_4^2}{16\pi^2} m_{1/2}\,, \tag{2.11}$$

where $m_{1/2}$ is the gaugino mass. The precise value of the counterterms is calculable if we fully specify the D-dimensional theory. However, we will see that the RG running of the soft masses from μ_c to the weak scale gives large additive contributions to the visible soft masses, and the final results are rather insensitive to the precise value of the counterterm. We will therefore be content with the simple estimate above.

Note that the soft terms arising from contact terms are larger than the anomaly-mediated contributions, which give

$$\Delta m_\lambda \sim \frac{g_4^2}{16\pi^2} \frac{F_{\rm hid}}{M_4}, \qquad \Delta m^2_{H_u}, \Delta m^2_{H_d} \sim \left(\frac{g_4^2}{16\pi^2} \frac{F_{\rm hid}}{M_4} \right)^2, \tag{2.12}$$

where $M_4 \gtrsim M$ is the 4-dimensional Planck scale [3, 5]. The other soft terms also get contributions larger than their anomaly-mediated values from the RG, as discussed above. Therefore, we can neglect the anomaly-mediated contribution in this class of models.

The higher-dimensional origin of these theories can also solve the 'SUSY CP problem' [18]. This problem arises from the fact that the phases in the SUSY breaking terms must be much less than 1, otherwise they give rise to electron and neutron electric dipole moments in conflict with experimental bounds. This is a naturalness problem because CP is (apparently) maximally violated in the CKM matrix, and it must be explained why it is not violated in all terms. In the present model, CP-violating phases can appear in μ, B, and $m_{1/2}$, generated from higher-dimension operators in eq. (2.6). The phases in μ and B can be rotated away using a combination of $\mathrm{U}(1)_{\rm PQ}$ and $\mathrm{U}(1)_R$ transformations, leaving a single phase in $m_{1/2}$. This phase can vanish naturally in the present model if CP is violated only by terms in the lagrangean localized on the visible brane. This is a natural assumption because loop effects do not generate local CP-violating terms in the bulk or the hidden brane. This situation can arise (for example) if CP is broken spontaneously by fields localized on the visible brane. (In order to avoid a large neutron electric dipole moment, we must also assume that the effects of the QCD vacuum angle are suppressed [19].)

There are many other aspects of the higher-dimensional theory that we could discuss, but the basic features of the scenario depend only on the qualitative feature that the visible and hidden sectors are spatially separated. A complete specification of the higher-dimensional model would have to take into account the fact that there

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are more supersymmetries in higher dimensions. This may be broken spontaneously or explicitly (e.g. by an orbifold), and couplings between bulk and boundary fields must be consistent with SUSY. An explicit example with 5 dimensions compactified on a $S^1/\mathbb{Z}_2$ orbifold is easily constructed [12, 9]. Another important feature of the higher-dimensional theory is the stabilization of the extra dimensions. Stabilization mechanisms that are appropriate for the scenario we are considering are discussed in refs. [20, 6]. We conclude that there is no obstacle to constructing realistic effective field theory models of the type outlined here. The question of whether a model of this type can be derived from a more fundamental theory such as string theory is left for future work.

3. Phenomenology

We now turn to the phenomenology of these models. We have seen that the SUSY breaking parameters in the effective 4-dimensional theory are determined at the compactification scale $\mu_c \sim 1/r$. We have also seen that μ_c is one to two orders of magnitude below the fundamental scale M, which is most naturally taken to be close to the string scale. Therefore, we expect μ_c to be close to the unification scale $M_{\rm GUT} \sim 2 \times 10^{16}\,\text{GeV}$. We therefore identify μ_c and $M_{\rm GUT}$ in making our estimates.

We will further assume that the theory is embedded in a grand-unified theory (GUT) at the scale $M_{\rm GUT}$, as suggested by the success of gauge coupling unification in the MSSM. We therefore consider the following SUSY breaking parameters renormalized at $M_{\rm GUT}$:

$$\begin{aligned} \text{Gaugino masses:}\quad & M_1 = M_2 = M_3 = m_{1/2}, \\ \text{Higgs masses:}\quad & m^2_{H_u}, m^2_{H_d} \sim m^2_{1/2}, \quad \mu, B \sim m_{1/2}, \\ \text{Squark and slepton masses:}\quad & m^2 \sim \frac{m^2_{1/2}}{16\pi^2}, \\ A \text{ terms:}\quad & A \sim \frac{m_{1/2}}{16\pi^2}. \end{aligned} \tag{3.1}$$

We have argued above that these conditions can emerge naturally in this scenario for $D = 5$ or 6. If we neglect the small loop-suppressed parameters, the model is defined by the 6 parameters $m_{1/2}$, $m^2_{H_u}$, $m^2_{H_d}$, μ, B, and y_t renormalized at $M_{\rm GUT}$. (We do not consider large $\tan\beta$ solutions, so we neglect all other Yukawa couplings.) The value of y_t at the weak scale fixes $\tan\beta$ from the observed value of the top quark. The requirement that electroweak symmetry breaks with the correct value of M_Z and $\tan\beta$ then fixes two more parameters. We see that we are left with essentially 4 parameters.

An important issue when analyzing the spectrum at the weak scale is the radiative corrections to the lightest neutral Higgs mass [21]. The largest effect can be viewed as a top loop contribution to an effective quartic term in the effective potential below the stop mass [22]. We include an estimate of this effect by adding the term

$$\Delta V_H = \left(\frac{3y_t^4}{8\pi^2} \ln \frac{m_{\tilde{t}}}{m_t} \right) (H_u^\dagger H_u)^2 \tag{3.2}$$

to the Higgs potential.

We evolve the 1-loop RG equations from the scale $M_{\rm GUT} = 2 \times 10^{16}\,{\rm GeV}$ down to the weak scale $\mu_{\rm W} = 500\,{\rm GeV}$, using $\alpha_{\rm GUT} = 1/(24.3)$. We use input values of $m_{1/2}$, $m_{H_u}^2$, $m_{H_d}^2$, and y_t at $M_{\rm GUT}$ and determine μ and B by imposing electroweak symmetry breaking. The value of the top quark mass is used to fix $\tan\beta$; we use $m_t(\mu_{\rm W}) = 165\,{\rm GeV}$, which includes 1-loop QCD corrections. We minimize the Higgs potential including the term eq. (3.2) with $m_{\tilde{t}}$ taken to be the heaviest of the stop mass eigenstates, and y_t renormalized at $\mu_{\rm W}$. These approximations could be refined, but they will suffice to illustrate the main features of the spectrum of this class of models.

Some parameter choices that give rise to realistic spectra are given in table 3. We find that the dependence on the overall scale of the initial SUSY breaking masses is what would be expected: the superpartners become heavier, and the amount of fine-tuning required to achieve electroweak symmetry breaking increases (see below). The right-handed sleptons get an important positive contribution from a hypercharge Fayet-Iliopoulos term if $m_{H_d}^2 > m_{H_u}^2$ at the GUT scale. This distinguishes this model from 'no-scale' models. This is illustrated in the second and third parameter points in table 3. For $m_{H_d}^2 > m_{H_u}^2$, we easily obtain spectra where the LSP is a neutralino. The value of y_t mainly influences the value of $\tan\beta$, which is important because the lightest Higgs boson is light for small $\tan\beta$. We also find that $\tan\beta \gtrsim 2.5$ is preferred in order to have a sufficiently large mass for the lightest neutral Higgs.

An important feature of these results is the amount of fine-tuning required to achieve electroweak symmetry breaking. We define the fractional sensitivity to a parameter c (a coupling renormalized at $M_{\rm GUT}$) to be [23]

$$\text{sensitivity} = \frac{c}{v} \frac{\partial v}{\partial c}\,, \tag{3.3}$$

where v is the Higgs VEV and the derivative is taken with all other couplings at the GUT scale held fixed. The largest sensitivity is to $m_{1/2}$ and μ, and the values of the sensitivity parameter are given in table 3. We see that the sensitivity increases strongly as the superpartner masses are increased. Note that even for parameters where the superpartner masses are close to the experimental limits, the sensitivity parameter is large ($\gtrsim 20$). However, it is argued by Anderson and Castaño in ref. [24]

		Point 1	Point 2	Point 3
inputs:	$m_{1/2}$	200	400	400
	$m^2_{H_u}$	$(200)^2$	$(400)^2$	$(400)^2$
	$m^2_{H_d}$	$(300)^2$	$(600)^2$	$(400)^2$
	μ	370	755	725
	B	315	635	510
	y_t	0.8	0.8	0.8
neutralinos:	$m_{\chi^0_1}$	78	165	165
	$m_{\chi^0_2}$	140	315	315
	$m_{\chi^0_3}$	320	650	630
	$m_{\chi^0_4}$	360	670	650
charginos:	$m_{\chi^\pm_1}$	140	315	315
	$m_{\chi^\pm_2}$	350	670	645
Higgs:	$\tan\beta$	2.5	2.5	2.5
	m_{h^0}	90	100	100
	m_{H^0}	490	995	860
	m_A	490	1000	860
	$m_{H^\pm}$	495	1000	860
sleptons:	$m_{\tilde{e}_R}$	105	200	160
	$m_{\tilde{e}_L}$	140	275	285
	$m_{\tilde{\nu}_L}$	125	265	280
stops:	$m_{\tilde{t}_1}$	350	685	690
	$m_{\tilde{t}_2}$	470	875	875
other squarks:	$m_{\tilde{u}_L}$	470	945	945
	$m_{\tilde{u}_R}$	450	905	910
	$m_{\tilde{d}_L}$	475	950	945
	$m_{\tilde{d}_R}$	455	910	905
gluino:	M_3	520	1000	1050
sensitivity:	$m_{1/2}$	16	50	50
	μ	19	78	78

Table 3: Sample points in parameter space. All masses are in GeV. In the first two points, the LSP is mostly Bino, while in the third it is a right-handed slepton. The sensitivity parameter is defined in the main text.

that sensitivity does not capture the idea of fine-tuning: the theory is fine-tuned only if the physical quantities significantly more sensitive than *a priori* allowed choices of parameters. From this point of view, the fine-tuning of points with low superpartner masses is much less severe, and naturalness clearly favors regions of parameters with light superpartner masses [24]. In particular, requiring that the naturalness parameter defined in ref. [24] be less than ~ 10 implies that the parameter $m_{1/2}$ should be less than $\sim 400\,\mathrm{GeV}$.

4. Conclusions

This model is the simplest supersymmetric theory in the literature that generates an acceptable spectrum for the superpartners while explaining the absence of non-standard flavor-changing processes and electric dipole moments. It is highly predictive, with squark and slepton masses qualitatively similar to those of 'no-scale' supergravity models. The nonuniversality of the up- and down-type Higgs masses at the GUT scale can distinguish this theory from 'no-scale' supergravity—the expected difference between the up and down type Higgs masses generates a hypercharge Fayet-Iliopoulos term which affects the slepton mass spectrum. The right-handed sleptons and the lightest neutralino are significantly lighter than the other superpartners, and obtaining natural electroweak symmetry breaking requires that these be lighter than roughly 200 GeV.

While this work was being completed, we received ref. [13], which considers very similar ideas.

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