NO-SCALE SUPERSYMMETRIC STANDARD MODEL

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We propose a class of supergravity models coupled to matter in which the scales of supersymmetry breaking and of weak gauge symmetry breaking are *both* fixed by dimensional transmutation, *not* put in by hand. The models have a flat potential with zero cosmological constant before the evaluation of weak radiative corrections which determine $m_{3/2}, m_W = \exp[-O(1)/\alpha_t] m_P : \alpha_t = O(\alpha)$. These models are consistent with all particle physics and cosmological constraints for top quark masses in the range 30 GeV $< m_t <$ 100 GeV.

The gauge hierarchy problem [1] has two aspects: the fixing of $m_W \sim O(10^{-17}) m_P$, and keeping it there despite the destabilizing influence of radiative corrections. Supersymmetry (SUSY) [2] provides a solution [3] to the second, technical aspect of this problem, since no-renormalization theorems [4] ensure that the gauge hierarchy is stable against radiative corrections. However, SUSY does not by itself explain the origin of the weak interaction scale. Conventional phenomenology does of course require that the mass splittings between observable particles and their spartners induced by local SUSY breaking must be above about 20 GeV, and the stability of the gauge hierarchy requires that the induced SUSY breaking splittings be less than about 1 TeV. In most models [‡1] the weak gauge symmetry breaking scale and/or the local SUSY breaking scale are put in by hand. This is clearly unsatisfactory, and one would like both scales to be determined dynamically. One scenario [6,7] for the dynamical determination of the weak gauge symmetry breaking scale is dimensional transmutation à la Coleman–Weinberg [8]. This can be implemented in a spontaneously broken supergravity [9] model, yielding a dynamically determined weak scale which is *a priori* unrelated to the scale of local SUSY breaking. However, in previous incarnations [6,7] of this scenario the local SUSY breaking scale was still put in by hand.

In this paper we propose a new realization of the dimensional transmutation scenario in which *both* the weak gauge symmetry breaking *and* the local SUSY breaking scales are fixed dynamically. We work in the context of a class [10] of $N = 1$ supergravity models whose potential is absolutely flat, with zero cosmological constant and an undetermined scale of local SUSY breaking at the tree level [‡2]. We use this as a hidden sector [9] coupled to the simplest possible observable sector which is the minimal SUSY standard model with no additional chiral superfields. When the t quark is heavy enough, non-gravitational radiative corrections in this model drive [11,12] the (mass)^2 of one of the

[‡1] For recent reviews see ref. [5].

[‡2] We do not consider gravitational radiative corrections to this flat potential, but regard it as an effective potential after all such effects are taken into account [11]. It may well be that the tree-level flatness of a least some of these [10] potentials is in any case preserved by gravitational radiative corrections, as we discuss in a moment.

Higgs fields negative, triggering spontaneous weak gauge symmetry breaking at some scale of order

$$Q_0 = \exp[-O(1)/\alpha_t]\, m_P : \alpha_t \equiv g^2_{H\bar{t}t}/4\pi = O(\alpha)\,. \quad (1)$$

The Higgs and other scalar masses are all proportional to the gravitino mass, which is free to be determined dynamically in these models [10] with a flat potential. The gravitino mass must be non-zero for weak gauge symmetry breaking to occur, but on the other hand it cannot be larger than the dimensional transmutation scale at which the effective Higgs (mass)^2 becomes negative. We in fact find

$$m_{3/2} = O(1) \times m_W = O(1) \times Q_0\,, \quad (2)$$

with the exact ratios being model-dependent as we shall see later in this paper. Thus we have a class of models in which the weak gauge and SUSY breaking scale are interrelated and dynamically determined (1), (2) to be hierarchically small.

We assume the existence of a hidden sector z whose Kähler potential $\mathcal{G}$ has the general form [10]:

$$\mathcal{G} = -\tfrac{3}{2}\log[f(z) + f^\dagger(z^*)]^2\,. \quad (3)$$

This yields a flat vanishing scalar potential

$$V = 9\exp(\tfrac{4}{3}\mathcal{G})\mathcal{G}^{-1}_{zz^*}\partial_z\partial_{z^*}\exp(-\tfrac{1}{3}\mathcal{G})\,,$$

and an arbitrary scale of local SUSY breaking at the tree level

$$m_{3/2} = \exp(\tfrac{1}{2}\mathcal{G}) = |f(z) + f^\dagger(z^*)|^{-3/2}\,. \quad (4)$$

All such theories with one hidden complex field z are in fact equivalent up to field redefinitions: $z \to f(z)$, and their kinetic terms possess an SU(1,1) symmetry [10]. This SU(1,1) symmetry is sufficient to guarantee a flat potential. Thus if the SU(1,1) symmetry were preserved by gravitational radiative corrections it would make a flat potential technically "natural". Our hidden sector is in fact identical with the scalar sector in an $N = 4$ extended supergravity theory [13], suggesting an intriguing link between our phenomenological use of simple supergravity and a more fundamental role for extended supergravity. It is natural to add this SU(1,1) symmetry to the list of desirable features which might be extracted from eventual superunification in a fundamental $N = 8$ extended supergravity theory: local SU(8) → SU(5)? [14], $N = 8$ SUSY broken to $N = 1$ SUSY at m_P? [15], and now [16,17] global $E_{7(7)} \to$ SU(1,1)?

We combine [10,16,17] with $\mathcal{G}(z)$ an observable sector y^i which has canonical kinetic terms and a cubic superpotential $h(y^i)$ that could simply be that of the SUSY standard model:

$$\mathcal{G}(y^i) = \ni h(y^i) + h^\dagger(y_i^*)\,. \quad (5)$$

The observable and hidden sectors should be combined in such a way that avoids Yukawa couplings which vary with $m_{3/2}$. This is done in ref. [17]. The potential for the matter fields becomes *at the tree level*

$$V(y^i, y_i^*) = \sum_i |\partial g/\partial y^i + m_{3/2} y_i^*|^2 + (A-3)[g(y^i) + g^\dagger(y_i^*)] + \tfrac{1}{2}D^\alpha D^\alpha\,, \quad (6)$$

where $g(y^i)$ is the rescaled superpotential. In a simple case [10] the parameter $A = 3$, but other values are possible [17]. The gravitino mass scale is as yet undetermined. The matter fields y could be invariant with respect to the conjectured SU(1,1) symmetry. This means that a non-trivial [9] kinetic term $f_{ab}(z,y)G^aG^b$ for the gauge fields G^a can only depend on the matter fields y: only $f_{ab}(z,y) = f_{ab}(y)$ is allowed by SU(1,1). This would suggest that $\xi = m_{\tilde{\gamma}}/m_{3/2} \simeq 0$, though a non-zero but presumably small value of ξ could be generated either by GUT y fields in f_{ab}, or by radiative corrections.

Focussing on the neutral Higgs fields in the SUSY standard model: $H_1 \leftrightarrow m_{d,s,b,e,\mu,\tau}$ and $H_2 \leftrightarrow m_{u,c,t}$ we get the potential [18]

$$V(H_1, H_2) = \tfrac{1}{8}(g_2^2 + g'^2)(|H_1|^2 - |H_2|^2) + m_1^2|H_1|^2 + m_2^2|H_2|^2 - 2m_3^2\mathrm{Re}(H_1^*H_2) \quad (7)$$

at the tree level. We recall that the SUSY breaking mass parameters m_i^2, which are proportional to $m_{3/2}^2$, are subject to radiative corrections which can be summed using the renormalization group [19], and can trigger weak gauge symmetry breaking if m_t is large enough [6,7,11,12,19]. If we adopt the minimal subtraction renormalization prescription then the one-loop radiative corrections $\delta V(H_1, H_2; Q)$ to the basic potential $V(H_1, H_2; Q)$ (7) take the form

$$\delta V(H_1, H_2; Q) = \frac{1}{64\pi^2}\sum_J (-1)^{2J}(2J+1)m_J^4[\ln(m_J^2/Q^2) - \tfrac{3}{2}]\,. \quad (8)$$

Weak gauge symmetry breaking takes place when $m_1^2 + m_2^2 > 2m_3^2$ and

$$m_1^2 m_2^2 < m_3^4 \qquad (9)$$

in which case

$$\langle 0|H_1|0\rangle = 2^{-1/2} v \cos\theta\,, \quad \langle 0|H_2|0\rangle = 2^{-1/2} v \sin\theta\,, \qquad (10)$$

where

$$v^2 = -8(m_1^2 \cos^2\theta - m_2^2 \sin^2\theta)/(g_2^2 + g'^2)\cos 2\theta\,,$$

$$\cos^2 2\theta = [(m_1^2 + m_2^2)^2 - 4m_3^4]/(m_1^2 + m_2^2)^2\,. \qquad (11)$$

The minimum value of $V(H_1, H_2; Q)$ is

$$V_{\min}(m_{3/2}; Q) = [-1/2(g_2^2 + g'^2)] \times [(m_1^2 - m_2^2) - (m_1^2 + m_2^2)|\cos 2\theta|]^2\,, \qquad (12)$$

where we have note explicitly that $V_{\min}$ depends on the value of the local SUSY breaking parameter $m_{3/2}$ which sets the overall scale for the m_i, and on the renormalization scale Q. What choice of Q makes the expression (12) the best approximation to the absolute minimum of the potential? We see from eqs. (8), (10) and (11) that there is some choice of $Q = O(m_{3/2})$ for which the one-loop radiative correction $\delta V(8)$ vanishes. We therefore evaluate (12) with the m_i renormalized at this $Q = O(m_{3/2})$, and see that $V_{\min}$ depends on $m_{3/2}$ both explicitly (since the m_i^2 are proportional to $m_{3/2}^2$ at the tree level) but also implicitly through the logarithmic renormalization group variation of the m_i^2. Let us write

$$t \equiv \ln(m_{3/2}^2/Q_0^2)\,, \qquad (13)$$

where Q_0 is the scale, determined by dimensional transmutation, at which the gauge symmetry breaking condition (9) is first satisfied. Then $V_{\min}$ takes the general form

$$V_{\min}(t) = e^{2t} F(t)\,, \qquad (14)$$

where $F(t)$ is a smooth and well-behaved function which vanishes for some value of $t = O(1)[m_{3/2} = O(Q_0)]$. The prefactor e^{2t} vanishes when $t \to -\infty$ ($m_{3/2} \to 0$) and $F(t)$ is negative (12) for intermediate values of t. If we now minimize $V_{\min}(t)$ (14) with respect to t, we find

$$F' + 2F = 0\,. \qquad (15)$$

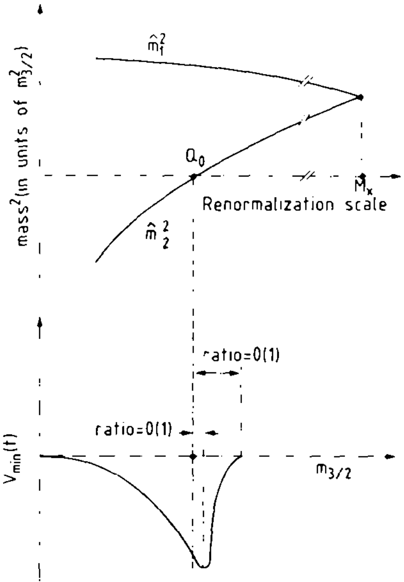

Fig. 1. Sketch of the variation of SUSY-breaking mass parameters m_i^2 with the renormalization scall Q. The Higgs (mass)^2 $m_2^2 = 0$ at a scale Q_0, which determines the dynamically preferred value of $m_{3/2}$, as seen in the bottom half of the figure.

The renormalization group tells us that $F' = O(\alpha)$, so the leading order version of the condition (15) is just $F = O(\alpha)$. This means that the absolute minimum of the potential is obtained when t is close to zero, i.e.

$$m_{3/2} = O(Q_0)\,, \qquad (16)$$

as illustrated in fig. 1. Thus the scale $m_s \equiv (m_{3/2} m_P)^{1/2}$ of local SUSY breaking is dynamically determined in terms of the dimensional transmutation scale (16).

More analysis is needed for the evaluation of the numerical ratio between $m_{3/2}$ and Q_0 (16). One simplifying feature is the observation that since $F(t) = O(\alpha)$ at the minimum, then $v^2 = O(m_{3/2}^2)$ and hence H field-dependent contributions to the particle masses included in the radiative corrections (8) can be neglected in leading order in α. Here we just present results for the idealized limit in which the gaugino mass $m_{\tilde{V}} = \xi m_{3/2} \to 0$ and also $m_3 \to 0$. In this case the gauge symmetry breaking condition (9) becomes

$$m_2^2 < 0\,, \qquad (9')$$

while

$$\langle 0|H_1|0\rangle \to 0\,,$$

$$(\langle 0|H_2|0\rangle)^2 \to -4m_2^2/(g_2^2 + g'^2) = \tfrac{1}{2}v^2 \qquad (10')$$

and

$$V_{\min}(m_{3/2}, Q) = -2m_2^4/(g_2^2 + g'^2)\,. \qquad (12')$$

The relevant dimensional transmutation scale Q_0 is now the scale at which $m_2^2(Q_0) = 0$. The previous analysis tells us that the absolute minimum of the potential occurs in the neighbourhood of this zero. Therefore we can approximate

$$m_2^2(Q^2) \approx C m_{3/2}^2 \alpha \ln(Q^2/Q_0^2)\,, \qquad (17a)$$

where C is a constant [7] of order unity:

$$C = (3/4\pi)(\alpha_t/\alpha)[(m_{\tilde{q}_3}^2 + m_{\tilde{u}_3}^2)/m_{3/2}^2 + A^2] + \ldots, \qquad (17b)$$

where $\alpha_t \equiv g_{H\bar{t}t}^2/4\pi$ is the t quark Yukawa coupling, $m_{\tilde{q}_3}$ and $m_{\tilde{u}_3}$ are the SUSY breaking mass parameters for the third-generation left-handed quark doublet and right-handed charge + 2/3 quark respectively, with these and A all evaluated at the renormalization scale Q_0. Substituting eq. (17) into (10′) we see that

$$v^2 = [8C/(g_2^2 + g'^2)]\, m_{3/2}^2 \alpha \ln(Q_0^2/Q^2)$$
$$\Rightarrow m_W^2 = \tfrac{1}{4} g_2^2 v^2 = 2\cos^2\theta_W m_{3/2}^2 \alpha \ln(Q_0^2/Q^2) \qquad (18)$$

and likewise $m_{\tilde{q}}^2, m_{\tilde{\ell}}^2 = O(\alpha) m_{3/2}^2$. Therefore, as already remarked, the H-dependent terms in eq. (8) can be neglected when we work to leading order in α. Using the results of ref. [7] we deduce that when $\xi, m_3^2 = 0$ and $m_2^2 = O(\alpha)$ (17a):

$$m_1^2 = m_{\tilde{d}_{1,2,3}}^2 = m_{\tilde{\ell}_{1,2,3}}^2 = m_{\tilde{u}_{1,2}}^2 = m_{\tilde{q}_{1,2}}^2 = m_{3/2}^2\,,$$
$$m_{\tilde{u}_3}^2 = \tfrac{1}{2} m_{3/2}^2\,,$$
$$m_{\tilde{q}_3}^2 = \tfrac{2}{3} m_{3/2}^2\,, \qquad (19)$$

and we compute that δV (8) vanishes for Q:

$$\ln(m_{3/2}^2/Q^2) = (\tfrac{3}{2} + \tfrac{9}{117}\ln 3 - \tfrac{8}{117}\ln 2) + O(\alpha)$$
$$\approx 1.537\,. \qquad (20)$$

Using this value of Q in eq. (17a), the expression (12′) for the minimum of the potential becomes

$$V_{\min}(t) \propto -e^{2t}(t - 1.537)^2\,. \qquad (14')$$

The condition (15) for the minimization with respect to $m_{3/2}$ becomes

$$t = 0.537\,. \qquad (15')$$

We are therefore able to relate $m_{3/2}$ directly to the dimensional transmutation scale:

$$m_{3/2}^2 = Q_0^2 \times e^{0.537}\,, \qquad (21)$$

while eq. (18) relates m_W to $m_{3/2}$. We therefore know what the dimensional transmutation scale must be in this model:

$$Q_0 = m_W \{2\pi/3[1 + A^2(Q_0)]\cos^2\theta_W \alpha_t\}^{1/2}$$
$$\times \exp(-0.268)\,. \qquad (22)$$

It is easy to determine in a given model what value of the t quark Yukawa coupling α_t will give the right value of Q_0. We assume for definiteness that the SUSY standard model is embedded in a GUT at some large scale m_X. Then there is a simple analytic expression [7] for $\alpha_t(Q_0)$ in the limiting case $\xi = m_{\tilde{V}}/m_{3/2} \to 0, m_3 \to 0$:

$$\alpha_t(Q_0) = \{r^2(Q_0)/[1 + 6r^2(Q_0)]\}$$
$$\times [\tfrac{16}{3}\alpha_3(Q_0) + 3\alpha_2(Q_0) + \tfrac{13}{15}\alpha_1(Q_0)]$$
$$\times \{1 - [\alpha_X/\alpha_3(Q_0)]^{16/9}[\alpha_X/\alpha_2(Q_0)]^{-2}$$
$$\times [\alpha_X/\alpha_1(Q_0)]^{-13/99}\}^{-1}\,, \qquad (23a)$$

where if $A = 3$ initially [10]

$$3r^2(Q_0)\{1 + \{3/[1 + 6r^2(Q_0)]\}^2\} = 1\,, \qquad (23b)$$

and α_X is the gauge coupling strenght at the grand unification scale, which yields

$$m_t \approx 110\ \text{GeV}. \qquad (24)$$

To see more explicitly the way in which dimensional transmutation determines Q_0 and hence the weak interaction scale via (22), we note that in the limit $m_3^2 \to 0, \xi \to 0$:

$$Q_0 = m_X \exp[-2\pi/3\bar{\alpha}_t(1 + \bar{A}^2) + \ldots]\,, \qquad (25a)$$

where $\bar{\alpha}_t$ and $\bar{A}$ are values of α_t and A at some intermediate scale $\bar{Q}$: $Q_0 < \bar{Q} < m_X$. The form (25a) parallels the corresponding relation between Λ_{QCD} and m_X in a SUSY GUT:

$$\Lambda_{QCD} = m_X \exp(-2\pi\sin^2\theta_W/3\alpha_{em} + \ldots)\,. \qquad (25b)$$

We see from (25) that for $\bar{A} = O(2)$ and $\sin^2\theta_W = O(0.2)$ it is no accident that Q_0 and Λ_{QCD} are similar in magnitude if $m_t \simeq m_W$.

Fig. 2 illustrates the values of m_t for $A = 3$ and general non-zero values of ξ and $\hat{m}_4 \equiv m_3^2/2m_{3/2}^2$. (Radi-

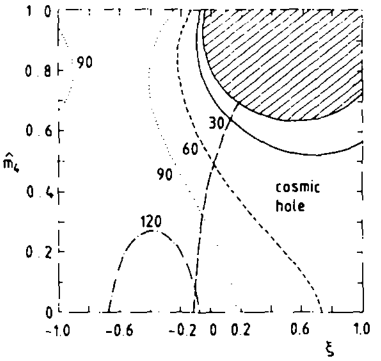

Fig. 2. Values of m_t (in GeV) for general values of $\hat{m}_4 \equiv m_3^2/2m_{3/2}^2$ and $\xi = m_{\tilde{\gamma}}/m_{3/2}$. There is no solution with $m_t > 20$ GeV in the shaded region, while our present vacuum is unstable in the cosmic hole region to the right of the broken line.

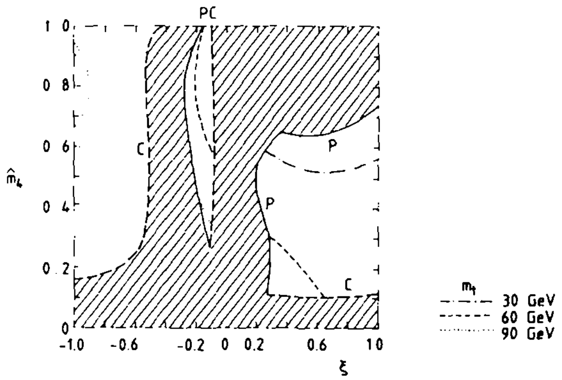

Fig. 3. Contours of m_t as in fig. 2, with the shaded domain indicating values of $\hat{m}_4$ and ξ disallowed because of the absence of a charged sparticle with mass less than 20 GeV (P), and/or because of an excessive cosmological density of the lightest neutral sparticle (C).

ative corrections in a GUT are able to generate $m_3 = O(m_{3/2})$.) The shaded region in the top right of fig. 2 corresponds to unacceptably small values of $m_t < 20$ GeV, while to the right of the broken line the desired $SU(2)_L \times U(1)_Y \to U(1)_{em}$ minimum is not the lowest, and tunneling into a lower minimum is possible [20]. This is not necessarily catastrophic [6], since the tunnelling time may be very long. We have evaluated the action in the thin-wall approximation for tunnelling into the $\langle 0|\tilde{t}\ |0\rangle \neq 0$ vacuum, and find it is always larger than 60. The persistence of our present vacuum to the ripe old age of 10^{10} years requires an action above 400. However, it has been observed [21] in a similar model that the thin-wall approximation underestimates the true action by a factor O(10) when the true action ≈ 400. Therefore we suspect that all values of ξ and $\hat{m}_4$ are acceptable from this cosmological point of view, though one might feel queasy in the cosmic hole region where our present vacuum is mortal. Fig. 3 shows which regions of the $\xi, \hat{m}_4$ plane are allowed by other phenomenological constraints, notably the absence of any charged supersymmetric particle with mass < 20 GeV denoted by P, the requirement [22] that the lightest neutral sparticle χ^0 be lighter than the lightest charged one $\chi^\pm$, and the requirement [22, 23] that the present cosmological density of this lightest neutral sparticle be less than 2×10^{-29} gm/cc, denoted by C. The requirement C excludes regions of small $|\xi|$ ($m_{\tilde{\gamma}} \to 0$), small $\hat{m}_4$ ($m_{\tilde{\nu}} \to 0$) and a band in the centre left of the figure where there is another light neutral sparticle (see graphs (b), (d) in refs. [22, 24]). The requirement P excludes may regions where different charged sparticles become too light. The requirement $m_{\chi^0} < m_{\chi^\pm}$ is almost automatic once P and C are satisfied. We see three allowed domains in fig. 3. The regions at large $|\xi|$ are disfavoured by our previous argument that SU(1,1) favours small ξ, while the right-hand region also has an unstable vacuum. We therefore prefer the thin central region of fig. 3, in which 33 GeV $< m_t <$ 103 GeV, and table 1 lists the complete sparticle spectrum for representative values of $|\xi|$ and m_3^2 corresponding to $m_t \approx 33$, 60 and 90 GeV. "Zen" decays [24, 25] of the $W^\pm$ into pairs of light gaugino/higgsino mixtures $\chi^\pm + \chi^0$ are allowed in all the preferred region of fig. 3, while $W^\pm \to \tilde{\ell}^\pm + \tilde{\nu}$ decays [26] are always kinematically impossible. Near the top of the preferred region in fig. 3 there is a neutral Higgs boson light enough to be detected in Υ decays [7]: we predict such a light Higgs if $m_t < O(40)$ GeV. There is also a domain near the bottom of the preferred region where the lighter stop squark is lighter than the t quark: $m_{\tilde{t}_1} < m_t$ [27].

There is less freedom in this class of models than in the traditional [6,7,11] models with gauge symmetry breaking driven by a heavy t quark. In our class of models small values of $\xi = m_{\tilde{\gamma}}/m_{3/2}$ are preferred, and the scale of $m_{3/2}$ is closely related to m_W. It is possible to avoid $A = 3$ [10] in more complicated models [17]. The observant reader will have noticed that, having started with a zero cosmological constant, radiative corrections have now regenerated $\Lambda = O(\alpha m_{3/2}^4)$. We expect that a shift in the cosmological constant $\delta\Lambda$

Table 1
Characteristic mass spectra (in GeV).

	m_t	33	60	90
	$m_{3/2}$	72	132	225
	$(\hat{m}_4, \xi)$	(0.9, −0.084)	(0.8, −0.195)	(0.4, −0.13)
sleptons	$\tilde{e}$	73	134	228
	$\tilde{e}^c$	73	133	228
	$\tilde{\nu}$	71	130	221
squarks	$\tilde{u}, \tilde{c}$	72	142	232
	$\tilde{d}, \tilde{s}$	75	146	239
	$\tilde{u}^c, \tilde{c}^c$	73	142	233
	$\tilde{d}^c, \tilde{s}^c$	74	144	235
	$\tilde{t}_1, \tilde{t}_2$	41, 100	93, 175	73, 278
	$\tilde{b}, \tilde{b}^c$	73, 74	136, 144	209, 235
higgses	$H^\pm$	156	236	291
	H^0	4.6, 134, 163	15, 221, 239	46, 290, 280
gauginos/	χ^0	3.2, 61, 66, 132	14, 39, 106, 164	16, 33, 98, 152
shiggses	$\chi^\pm$	51, 121	27, 154	25, 139
	$\tilde{g}$	15	65	74

$= O(\alpha m_{3/2}^4)$ would yield a shift in the gravitino (mass)^2 $\delta m_{3/2}^2 = O(\alpha m_{3/2}^4/m_P^2)$. This would be negligible by comparison with $m_{3/2}^2$ itself and hence have negligible impact on the scenario described in this paper.

These are the first models where *neither* the SUSY breaking scale *nor* the weak gauge symmetry breaking scale is put in by hand, but both are determined by radiative corrections. Thus these are "no-scale" SUSY standard models with all light mass parameters provided by dimensional transmutation. Clearly it would be interesting challenge to construct an analogous "no-scale" SUSY GUT where m_X is related to m_P in some way. This is probably not very difficult so we regard this SUSY standard model as an existence proof for theories of all the elementary particle interactions whose only intrinsic scale is the Planck mass.

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