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WEAK SYMMETRY BREAKING BY RADIATIVE CORRECTIONS IN BROKEN SUPERGRAVITY*

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ABSTRACT

Weak interaction gauge symmetry breaking can be generated by radiative corrections in a spontaneously broken supergravity theory, provided the top quark is heavy enough. In one class of such theories the weak Higgs vacuum expectation values are determined by dimensional transmutation à la Coleman-Weinberg, and may be considerably larger than the magnitudes of susy breaking mass parameters. In this scenario $m_t \geq 65\ GeV$, the supersymmetric partners of known particles may have masses $\ll m_W$, the mass of the lighter neutral scalar Higgs boson is determined by radiative corrections, and there is some variant of a light pseudoscalar axion. In contrast to conventional Coleman-Weinberg models, the weak phase transition is second order and there is no likelihood of excess entropy production.

Supersymmetry (susy) has recently attracted considerable phenomenological attention because[1] it can protect the weak interaction scale and preserve the hierarchy $m_W/m_P \ll 1$. However, susy does not by itself predict or explain the magnitude of m_W. Also, although the susy partners of many familiar particles must have masses $\leq O(1)\ TeV$ if the hierarchy is to be maintained, the primordial susy breaking scale $\sqrt{d}$ could be much larger.[2,3] Thus, scenarios have been proposed in which the weak interaction scale is obtained from high order radiative corrections,[3] with symmetry breaking driven by a heavy top quark.[4,5] When $\sqrt{d} \geq O(10^{11})\ GeV$ it seems essential to consider the effects of local susy, since the gravitino mass $m_{3/2} = O(d/m_P) = O(m_W)$, and scalar fields acquire contributions m to their masses of $O(m_{3/2})$.[6] Some phenomenological supergravity models have been proposed in which weak gauge symmetry breaking is realized at the tree level.[7] However, it seemed[8] to us more natural to suppose that radiative corrections play an important role, possibly with a heavy t quark driving weak gauge symmetry breaking, as had been proposed earlier[3,4,5] in the context of global susy (see also ref. 9). Moreover, there emerged[10] difficulties with alternative models for weak symmetry breaking which employed light singlet chiral superfields. In the previous paper[8] we demonstrated the feasibility of a similar scenario in the context of local susy, without solving the full coupled set of renormalization group equations for the susy breaking parameters.

Conveniently enough, the full renormalization group equations for these parameters are available from a previous analysis[5] in the context of global susy. All that is necessary in order to arrive at an analogous broken supergravity model is to choose a somewhat different set of initial conditions for the susy breaking parameters.[11] These include gaugino masses M,[8] scalar boson masses m,[6] and trilinear scalar couplings λ.[7,12,13] One's guess might be that all of these parameters are $O(m_{3/2})$. However, it has been proposed[14] on the basis of a $U(n)$ symmetry among the chiral superfields respected by perturbative gravitational effects, that perhaps $M = O(\alpha/2\pi)m_{3/2}$.[13] We see no particular reason why such a symmetry should survive non-perturbative gravitational effects, and it is in any case broken by Yukawa couplings which may be large for the top quark. Therefore we prefer to retain $\hat{M} \equiv M/m = O(1)$. The initial value of the ratio $\hat{\lambda} \equiv \lambda/m$ is related[7,12,13] to unknown parameters of a hidden sector of the theory, and is model-dependent but probably $O(1)$.

We prefer to keep an open mind about this sector of the theory, which may well not be a simple polynomial in a single unknown chiral superfield added on to the superpotential for known chiral superfields,[15] but may reflect some more complicated dynamics at scales $O(m_P)$. In addition to the mass parameters listed above, the low energy Higgs potential involving two Higgs superfields $H_{1,2}$ with susy breaking masses $m_{1,2}$ may also include a quadratic term H_1H_2 with coefficient $\mu \times O(m_{3/2})$ related to a quadratic term $\propto H_1H_2$ in the chiral superpotential. There is no *a priori* connection between the values of μ and of $m_{3/2}$, and if $\mu \ll m_W$ the physical Higgs spectrum contains an axion. Phenomenological model-builders search in the multi-dimensional space of the parameters m, $\hat{M}$, $\hat{\lambda}$, μ and the t quark Yukawa coupling h_t to the Higgs H_2 for outputs of the renormalization group equations in which m_2^2 has been driven negative by h_t, permitting the breakdown of $SU(2) \times U(1)$ to $U(1)_{em}$. Typically, for given choices of m, $\hat{M}$, $\hat{\lambda}$ and μ we find a range of values of h_t which give m_2^2 negative, corresponding to $m_t \geq O(M_W)$. Since m_2^2 varies quite rapidly as one approaches the strong interaction scale, different negative values of m_2^2 are attained at the price of modest variations in h_t and hence m_t.

Instead of reporting on a general survey[16] of this parameter space, we have chosen to formulate plausible hypotheses which diminish its dimensionality and constrain the theory in an interesting way. Since μ has no definite reason to be $O(m_W)$, and could well be much less, perhaps $O((\alpha/\pi)^n)m_W$ or $O(m_W^2/m_X)$ or even zero, we consider the possibility

$$\mu = 0, \quad \textit{or at least} \quad \ll m_W \;. \tag{1}$$

In this case the weak gauge symmetry breaking occurs near a scale Q_0 where the linear combination $m_1^2 + m_2^2$ of Higgs mass^2 parameters vanishes. This scale Q_0 is independent of m as long as $m \ll Q_0$. Furthermore, for a given choice of $\hat{M}$ and $\hat{\lambda}$ there is a *unique* value of h_t and hence m_t which fixes Q_0 so as to give m_W correctly. This enables us to *predict* m_t as a function of $\hat{M}$ and $\hat{\lambda}$, and we find that for all plausible values of these parameters

$$m_t \geq 65 \; GeV \;. \tag{2}$$

In contrast to other models, in this scenario the unseen supersymmetric partners of known particles could be lurking arbitrarily close to the present experimental lower limits on their masses. In this scenario the weak interaction scale is divorced from the scalar and gravitino masses, since it is fixed by dimensional transmutation in the style of Coleman and E. Weinberg[17] The difference is that whereas in their case it was the logarithmic evolution of a *quartic* Higgs coupling which determined the weak interaction scale, in our susy case it is the logarithmic evolution of a *quadratic* Higgs coupling. As in the Coleman-Weinberg analysis, we have a light neutral Higgs scalar whose mass is determined by radiative corrections, and we also have the pseudoscalar axion mentioned earlier. We assume that this axion could ultimately be made phenomenologically acceptable, perhaps by becoming a new improved invisible axion in a GUT[18,19] or perhaps by μ being sufficiently large ($\geq O(1)MeV$) to push the axion mass $m_a = O(\mu m)^{1/2}$

above the experimental lower limit of 350 MeV from $K \to \pi + a$ decay. It is interesting to speculate that the initial stage of GUT symmetry breaking could also be driven by radiative corrections, in which case one might hope to understand why $m_W/m_X \ll m_X/m_P \ll 1$ along the lines proposed in ref. 20. In this connection we make some remarks about the variation in couplings and mass parameters between m_P and m_X.

We assume there are no other light chiral superfields besides the Higgses $H_{1,2}$, the quarks and the leptons. Therefore the low energy potential for the neutral Higgses is[21]

$$V = \frac{g_2^2 + g'^2}{8}(|H_1|^2 - |H_2|^2)^2 + m_1^2|H_1|^2 + m_2^2|H_2|^2 - m_3^2(H_1H_2 + H_1^*H_2^*) \quad (3)$$

The quartic D-term allows the Higgses to leak to infinity unless[21]

$$m_1^2 + m_2^2 > 2m_3^2 \quad (4)$$

and there is $SU(2) \times U(1)$ breaking if[21]

$$m_3^4 > m_1^2 m_2^2 \quad (5)$$

with

$$\frac{v_1}{v_2} \equiv \frac{< 0|H_1|0 >}{< 0|H_2|0 >} = \cot\theta : \sin 2\theta = \frac{2m_3^2}{(m_1^2 + m_2^2)} \quad . \quad (6)$$

We assume that H_2 gives mass to the t quark $m_t = (1/\sqrt{2})h_t v_2$, and $h_t > h_b$ so that the renormalization group drives $m_2^2 < m_1^2$ at present energies, and we will be interested in what happens when $m_3^2 = O(\mu m) \to 0$. In leading order of the renormalization group equations the Higgs mass parameters m_i^2 in the effective potential depend (logarithmically) only on the corresponding $|H_i|^2$, and they are positive at large scales ensuring that condition (4) is obeyed. If $m_1^2 + m_2^2$ decreases to zero at some scale $|H_i| = Q_0$, this will determine the value of $v^2 \equiv v_1^2 + v_2^2$, while

$$\delta^2 \equiv v_2^2 - v_1^2 = \frac{2(m_1^2 - m_2^2)}{(g_2^2 + g'^2)} = \frac{4m_1^2(Q_0^2)}{(g_2^2 + g'^2)} \quad . \tag{7}$$

The combination $m_1^2 + m_2^2$ becomes negative at scales less than Q_0, resulting in the form of potential shown in fig. 2. If $m_{1,2}$ are much less than the dimensional transmutation scale Q_0 then equation (7) tells us that the absolute minimum of the potential is at

$$v_1^2 \approx v_2^2 \approx \frac{v^2}{2} \tag{8}$$

and

$$Q_0 = \sqrt{\frac{e}{2}}\, v \approx 290\ GeV\ . \tag{9}$$

To calculate the scale Q_0 at which $m_1^2 + m_2^2 = 0$ we need the leading order renormalization group equations of ref. 5 which are valid for $Q \gg M^2, m^2$. We have in their notation the initial conditions

$$m_3 = m_4 = m_5 = m_7 = m_9 = 0;\ m_6 = m_8 = m_{10} = \hat{\lambda}\, m \tag{10}$$

In the limit that $m = 0$ our initial conditions become a limiting case of those considered in ref. 5, with the only susy breaking in the initial conditions coming from $M \neq 0$. Neglecting all Yukawa couplings except those of the top quark, the relevant renormalization group equations are

$$Q\frac{dm_1^2}{dQ}=\frac{1}{(4\pi)^2}[-6g_2^2M_2^2-2g'^2M_1^2] \tag{11a}$$

$$Q\frac{dm_2^2}{dQ}=\frac{1}{(4\pi)^2}[-6g_2^2M_2^2-2g'^2M_1^2+6h_t^2(m_{q_3}^2+m_{p_3}^2+m_2^2+m_{10}^2)] \tag{11b}$$

$$Q\frac{dm_{10}}{dQ}=\frac{1}{(4\pi)^2}\Big[\frac{-32}{3}g_3^2M_3-6g_2^2M_2-\frac{26}{9}g'^2M_1+6h_t^2m_{10}\Big] \tag{11c}$$

$$Q\frac{dm_{q_3}^2}{dQ}=\frac{1}{(4\pi)^2}\Big[\frac{-32}{3}g_3^2M_3^2-6g_2^2M_1^2-\frac{2}{9}g'^2M_1^2+2h_t^2(m_{q_3}^2+m_{p_3}^2+m_2^2+m_{10}^2)\Big] \tag{11d}$$

$$Q\frac{dm_{p_3}^2}{dQ}=\frac{1}{(4\pi)^2}\Big[\frac{-32}{3}g_3^2M_3^2-\frac{32}{9}g'^2M_1^2+4h_t^2(m_{q_3}^2+m_{p_3}^2+m_2^2+m_{10}^2)\Big] \tag{11e}$$

$$Q\frac{dm_{n_3}^2}{dQ}=\frac{1}{(4\pi)^2}\Big[\frac{-32}{3}g_3^2M_3^2-\frac{8}{9}g'^2M_1^2\Big] \tag{11f}$$

for the susy breaking scalar mass parameters, and

$$Q\frac{dh_t}{dQ}=\frac{h_t}{(4\pi)^2}\Big[\frac{-16}{3}g_3^2-3g_2^2-\frac{13}{9}g'^2+6h_t^2\Big] \tag{12}$$

for the t quark Yukawa coupling. The gaugino masses are

$$M_{3,2}=\frac{g_{3,2}^2(Q^2)M}{g_{GUT}^2}\ ,\quad M_1=\frac{5}{3}\frac{g'^2(Q^2)M}{g_{GUT}^2} \tag{13}$$

while $g_{3,2}$ and g' evolve conventionally with Q.

We have integrated these renormalization group equations for different starting values of the ratios $\hat{M}$ and $\hat{\lambda}$, and located the corresponding values of m_t which yield a dimensional transmutation scale $Q_0=290\ GeV$. Vacuum stability conditions prefer[22] $\hat{\lambda}<3$, but this condition should be interpreted *cum grano salis.* It is applicable at scales $O(m_W)$ where $\hat{\lambda}$ is renormalized from its initial value in different ways for different trilinear couplings. Finite temperature effects in the early universe favour the conventional local minimum. Tunnelling into other minima is suppressed by $\exp(-O(1)/h^2)$ where h is the relevant Yukawa coupling. The false vacuum is more stable than the age of the universe except perhaps for transition to the minimum controlled by h_t. If $m_2^2 \ll m_{q3}^2, m_{p3}^2$ at scales $O\ (m_W)$, the absolute stability condition[22] on $\hat{\lambda}$ is modified to $\hat{\lambda}_t \equiv \lambda_t(m_W)/m_{\tilde{t}}(m_W) < 2$. This condition is obeyed if the initial $\hat{\lambda} \leq 2(1/2)$, as can be seen in the Table. Even if this condition is not obeyed, it is still possible that the lifetime of the false vacuum may be longer than the age of the universe for relevant values of h_t.

Our results for m_t are shown in fig. 1: they were determined by integrating the renormalization group equation for h_t down to a momentum scale $Q = m_t$. Note that we are not able to find solutions if $\hat{M} < 0.35$ for $\hat{\lambda} = 1$. Within the allowed range of $\hat{M}$ we find $m_t \geq 65\ GeV$ in the supersymmetric Coleman-Weinberg scenario for $\hat{\lambda} < 2\frac{1}{2}$. If m_t turns out to be $< 65\ GeV$, our scenario could still apply if our present vacuum is unstable, or if there is a fourth generation. In general, m_2^2 is evolving very rapidly at low Q, which means that the values of m_t needed are not much larger than the typical ranges found when we look for general solutions[5,16] to the inequalities (4,5) rather than looking specifically for $m_1^2 + m_2^2 \to 0$. In the general case we often find $v_1 \ll v_2 \approx v$, so that the same value of h_t gives m_t a factor $\sqrt{2}$ larger than in the dimensional transmutation case (8).

The rapid final stages of evolution of m_2^2 are driven by the increases in the t quark Yukawa coupling and more importantly in the squark masses which occur when $g_3^2/4\pi$ becomes large. Thus in the supersymmetric Coleman-Weinberg scenario the weak interaction scale is related to that of the strong interactions, while the absolute values of m and $m_{3/2}$ are not directly related to m_W. This contrasts with what usually happens in models of weak gauge symmetry breaking in supergravity models[7,8,9,16] where m_W is connected with m and $m_{3/2}$, but is not directly related to the strong interaction scale. In practice, phenomenology dictates that m must be large enough for all unobserved particles to have been able to escape detection, but it could be as low as 15 GeV in our scenario, thus offering the prospect of imminent detection of susy particles. The table shows values of the physical masses of these particles in units of m for selected representative values of the input parameters $\hat{M}$ and $\hat{\lambda}$. We see that the lightest spin-zero superpartners are the sleptons. For small M the lightest gaugino is approximately a photino $\tilde{\gamma}$ with mass

$$m_{\tilde{\gamma}} \approx \frac{g'^2 M_2 + g_2^2 M_1}{(g_2^2 + g'^2)} \approx \frac{8}{3} \frac{g_2^2 g'^2}{(g_2^2 + g'^2) g_{GUT}^2} M \approx 0.47 M \ . \tag{14}$$

This could be light enough to be pair-produced at PEP and PETRA, and the selectron mass could well be small enough for the cross-section for $e^+e^- \to \tilde{\gamma}\tilde{\gamma}\gamma$ to be detectably large at present energies.[23] Turning now to the physical Higgs bosons in this class of model,[21] the charged bosons $H^\pm$ and the heavier neutral scalar boson $H^{0\prime}$ acquire masses

$$m_{H^\pm} = m_{W^\pm}, \qquad m_{H^{0\prime}} = m_{Z^0} \tag{15}$$

at the tree level. The lighter extra scalar boson H^0 acquires

$$m_{H^0}^2 = \frac{1}{16\pi^2} \left\{ 6h_t^2 (m_{q_3}^2 + m_{p_3}^2 + m_2^2 + m_{10}^2) - 12 g^2 M_2^2 - 4 g'^2 M_1^2 \right\} \tag{16}$$

from radiative corrections. Values of m_{H^0} corresponding to typical values of the input parameters $\hat{M}$ and $\hat{\lambda}$ are also given in the table. Typically

$$m_{H^0} \approx \left(\frac{1}{4} \ to \ 2/3 \right) m \tag{17}$$

which is not much smaller than the slepton masses, as a result of the relatively large squark masses exhibited in the table and appearing in eq. (16). Finally, our spectrum contains a light neutral pseudoscalar axion state which must be exorcised in one of the ways discussed earlier. This can[19] be done in such a way as to avoid astrophysical and cosmological pitfalls. Our class of susy Coleman-Weinberg models also avoids the danger[24] of excess entropy generation during the weak phase transition, because as seen from fig. 2 the origin is an unstable extremum and there is a second order phase transition once the temperature falls below $\mathcal{O}(m)$.

Before closing we would like to add a few comments about the possibility of embedding this susy Coleman-Weinberg scenario in a GUT. One remark contains the initial values of the scalar masses that we have assumed. There is no good reason why the masses of $\underline{\overline{5}}$ and $\underline{10}$ matter fields $\overline{F}$ and T should be the same at the GUT breaking scale m_X, nor why the $\underline{5}$ and $\underline{\overline{5}}$ Higgs masses should be the same. Even if some symmetry fixed them to be equal at m_P, they would differ at M_X. We have evaluated this possible difference in the minimal $SU(5)$ GUT[1] and found that

$$1 \leq \hat{m}_T^2 \equiv \frac{m_T^2(m_X)}{m_{\overline{F}}^2(m_X)} \leq 1.5 \tag{18}$$

with $m_{H_1}^2 \approx m_{H_2}^2 \approx m_{\overline{F}}^2$. Figure 1 shows that variation in the range (18) does not have a substantial effect on the required t quark mass, though it can increase the physical masses of squarks and sleptons from the $\underline{10}$ representations of $SU(5)$, such as the $\tilde{e}_R$, $\tilde{\mu}_R$ and $\tilde{\tau}_R$.

It is enticing to speculate whether the grand unification scale m_X could also be determined by dimensional transmutation, thanks to some susy breaking scalar mass in the GUT sector being driven to zero at a scale $Q = \mathcal{O}(m_X)$. This would be a reincarnation of the double Coleman-Weinberg scenario of ref. 20, in which the "hierarchy of hierarchies" $m_W/m_X \ll m_X/m_P \ll 1$ was ascribed to the rapid evolution of the couplings of large GUT representations such as the $\underline{24}$ of Higgs in $SU(5)$ which gave a very large dimensional transmutation scale to the GUT breaking. This suggestion would now be applied to the susy breaking mass parameters instead of the quartic scalar couplings as illustrated in fig. 3. Unfortunately, such a scenario cannot be realized in the minimal susy GUT[25] where the lightness of the Weinberg-Salam Higgses and the heaviness of their colour triplet partners are enforced by the fine-tuning of two mass parameters in the superpotential. If one supplements the conventional minimal $SU(5)$ GUT with additional $\underline{40}$ and $\overline{\underline{40}}$ chiral superfields with a coupling ν to the adjoint $\underline{24}$ of Higgs, one can easily find plausible initial conditions at m_P which can drive m_{24}^2 to zero at scales $Q = O(10^{-3})m_P$, such as

$$\frac{g_5^2}{4\pi} = 0.19 \ ; \ \frac{\nu^2}{4\pi} = 0.004 \ ; \ m_{40}^2 = m_{\overline{40}}^2 = other\ m^2 \ ; \ M = O(2)\, m \quad . \quad (19)$$

It remains to find a cleverer model featuring such a supersymmetric hierarchy of hierarchies in which the Higgs doublet/triplet splitting problem is also solved.

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TABLE

Masses in models with the SU(2) × U(1) breaking

scale determined by radiative corrections.

	$\hat{M} = \hat{\lambda} = 1$	$M = 1$, $m = \lambda = 0$	$\hat{M} = .35$, $\hat{\lambda} = 1$	$\hat{M} = 1$, $\hat{\lambda} = 2.5$	$\hat{\lambda} = \hat{m}_T = \sqrt{3/2}$, $M = m_{\bar{F}}$
m_t	88	67	140	70	82
$\hat{m}_{H_o}$	.67	.46	.26	.66	.67
$\hat{m}_{\tilde{q}_3}$	2.7	2.6	1.1	2.7	2.7
$\hat{m}_{\tilde{q}_{1,2}}$	2.9	2.7	1.4	2.9	3.0
$\hat{m}_{\tilde{p}_3}$	2.3	2.4	.60	2.3	2.4
$\hat{m}_{\tilde{p}_{1,2}}$	2.8	2.6	1.4	2.8	2.9
$\hat{m}_{\tilde{n}_{1,2,3}}$	2.8	2.6	1.4	2.8	2.8
$\hat{m}_{\tilde{\ell}_{1,2,3}}$	1.2	.73	1.0	1.2	1.2
$\hat{m}_{\tilde{e}_{1,2,3}}$	1.1	.39	1.0	1.1	1.3
$\hat{\lambda}_t$	1.6	1.5	1.0	2.2	1.7

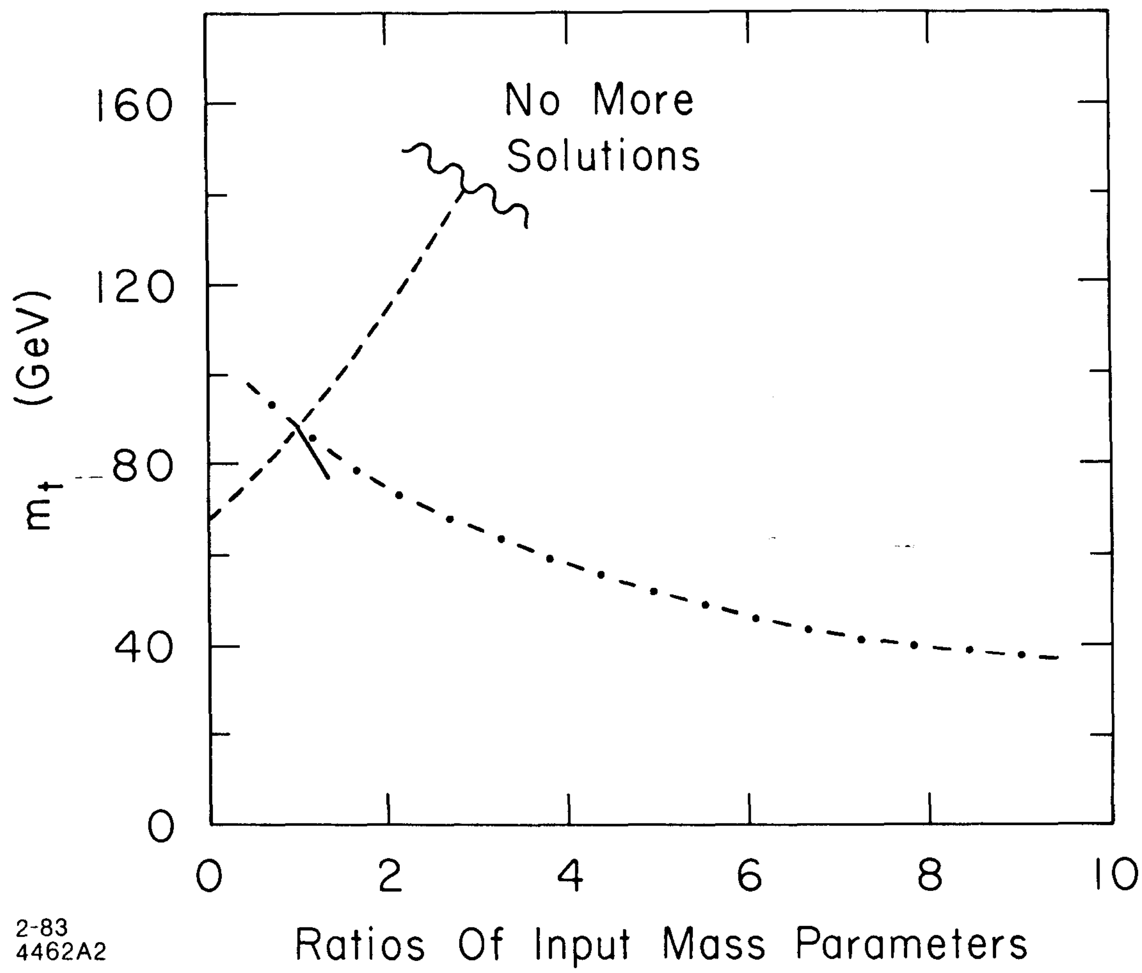
All masses denoted $\hat{m}_i$ are in units of the 5-plet scalar masses at the grand unification scale m_X, except that masses in the second column are in units of the gaugino mass at the scale m_X.

FIGURE CAPTIONS

1. Predictions of m_t corresponding to different values of the input mass ratios $\hat{M}^{-1} \equiv m/M$ (dashed line), $\hat{\lambda} \equiv m_{10}/m$ (dashed-dotted line) and $m_T/m_{\bar{F}}$ (solid line).

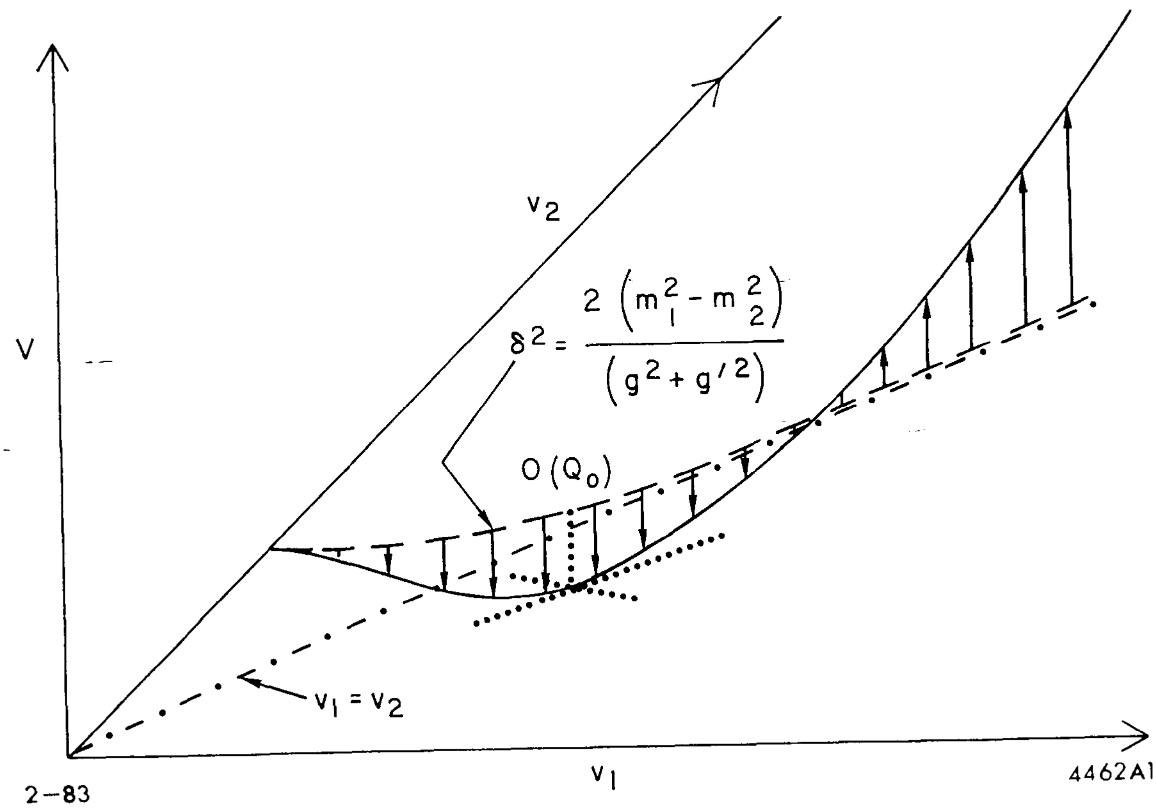
2. Form of potential in the dimensional transmutation scenario. The dashed line represents the curve of minima (7) in the (v_1, v_2) plane. The solid line represents the shape of the potential along this curve induced by the radiative corrections (11a) and (11b). The dotted lines show the location and depth of the absolute minimum of the potential at $\mathcal{O}(Q_0)$ where $m_1^2 + m_2^2 \simeq 0$. The extremum at $v_1 = 0, v_2 \neq 0$ is unstable since $m_1^2 + m_2^2 < 0$ at scales $\mathcal{O}(m) \ll Q_0$.

3. Qualitative features of the variation of susy breaking mass parameters in the "hierarchy of hierarchies" scenario. It may be possible to generate $m_X = \mathcal{O}(Q_1)$: $Q_0/m_P \ll Q_1/m_P \ll 1$.



No More Solutions
160
120
80
40
0
m_t (GeV)
0
2
4
6
8
10
2-83
4462A2
Ratios Of Input Mass Parameters

Fig. 1



V
v_2
$\delta^2 = \dfrac{2\left(m_1^2 - m_2^2\right)}{\left(g^2 + g'^2\right)}$
$O(Q_0)$
$v_1 = v_2$
v_1
2–83
4462A1

Fig. 2

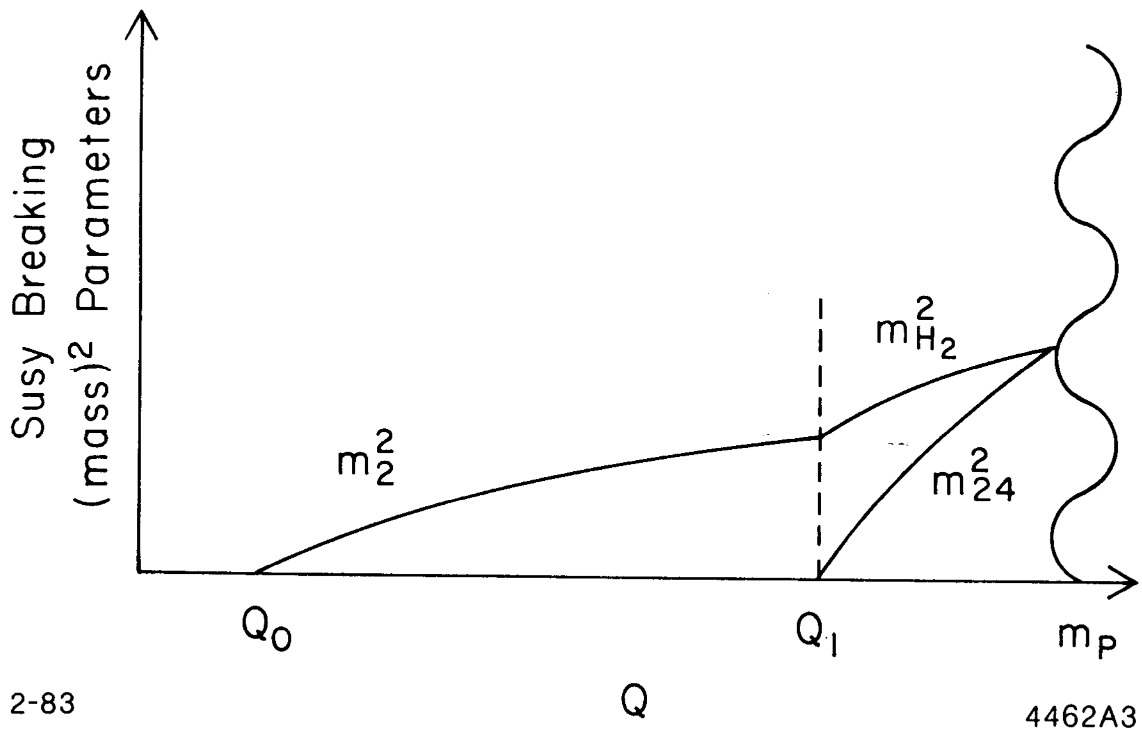

Susy Breaking
(mass)^2 Parameters
m_2^2
$m_{H_2}^2$
m_{24}^2
Q_0
Q_1
m_P
2-83
Q
4462A3

Fig. 3