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SUPERGRAVITY INDUCED RADIATIVE $SU(2) \times U(1)$ BREAKING

WITH LIGHT TOP QUARK AND STABLE MINIMUM

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A B S T R A C T

We present a class of models where the $SU(2)_L \times U(1)_Y$ breaking occurs radiatively through supergravity effects. Unlike similar models proposed up to now, our class of models : a) is free from a lower bound on the top quark mass (the only bound being the experimental one $m_t \geq 20$ GeV) ; b) enjoys a stable minimum ; and c) provides an experimentally accessible low energy particle spectrum.

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The introduction of supersymmetry into grand unified theories (GUTs) is, nowadays, the only realistic way of solving the gauge hierarchy problem ¹⁾. Non-renormalization theorems in spontaneously or softly broken supersymmetric theories guarantee the stability of the gauge hierarchy under radiative corrections while not shedding any light on the problem of doublet-triplet splitting at the tree level.

All successful global supersymmetric models have pointed towards a large scale of supersymmetry breaking $M_S \sim (M_W M_P)^{\frac{1}{2}} \approx 10^{10}$ GeV indicating thus that gravitational interactions should be of some relevance ²⁾. This led people to couple GUTs to $N = 1$ supergravity. Then, when local supersymmetry is spontaneously broken by the super-Higgs phenomenon ³⁾, the effective low energy theory becomes a global supersymmetric theory softly broken by the gravitino mass induced terms ⁴⁾. In this context also, the doublet-triplet hierarchy problem at the tree level has recently been solved using a non-renormalizable coupling in local supersymmetric GUTs ⁵⁾.

Lately great effort has been devoted towards the construction of local supersymmetric GUTs with a correct cosmological behaviour ⁶⁾. We prove in a separate publication ⁷⁾ that these cosmologically acceptable models are realistic in the sense that they lead to correct GUT predictions and phenomenologically acceptable behaviour.

Hereafter we shall concentrate on the $SU(2)_L \times U(1)_Y$ breaking. Two classes of mechanisms have been proposed so far.

a) Tree-level breaking, by using a light singlet coupled to the Higgs doublets.

However, this mechanism generates a mass term for doublets proportional to M_S destabilizing the hard-won gauge hierarchy and becoming thus useless for GUTs ⁸⁾.

b) Radiative breaking which is left as the only open possibility when

$SU(2)_L \times U(1)_Y$ is embedded in a grand unification group. The so far proposed models ⁹⁾⁻¹¹⁾ using this radiative mechanism require a large top Yukawa coupling and hence a large top quark mass, to drive the Higgs mass squared negative, resulting in $SU(2)_L \times U(1)_Y$ breaking. In this letter we will show that this is not a generic feature. Actually, we propose a general scheme of radiative breaking which allows for a light top quark.

The low energy potential for the neutral Higgses H_1 and H_2 is

$$V_{\text{tree}} = m_1^2 |H_1|^2 + m_2^2 |H_2|^2 - m_3^2 (H_1 H_2 + \text{h.c.}) + \frac{\hat{g}^2}{2} D^2 \quad (1)$$

where $D \equiv |H_1|^2 - |H_2|^2$ and $\hat{g}^2 \equiv (g_2^2 + g'^2)/4$ *).

The condition for the potential to be bounded from below is

$$m_1^2 + m_2^2 - 2m_3^2 > 0 \quad (2)$$

and the signal for $SU(2)_L \times U(1)_Y$ breaking is given by

$$\det \mathcal{M}^2 \equiv m_1^2 m_2^2 - m_3^4 < 0 \quad (3)$$

where $\mathcal{M}^2$ is the Higgs mass matrix squared.

Information coming from supergravity breaking imposes the initial condition $m_1^2 = m_2^2$ at scales $\sim M_X$ and this on account of the stability condition (2) requires $m_3^2 < m_{1,2}^2$ at M_X . This guarantees also, via inequality (3), that $SU(2)_L \times U(1)_Y$ remains unbroken at these high scales. However, taking into account the evolution of the parameters m_1^2 , m_2^2 and m_3^2 from M_X down to $\sim M_W$ by the renormalization group equations ¹²⁾, Eq. (3) can be satisfied at low energies so that symmetry breaking can occur.

In the literature, the case $m_3^2 = 0$ has been extensively studied ^{10), 11)}. In this particular case ($m_3^2 = 0$) an $SU(2)_L \times U(1)_Y$ breaking minimum starts to be developed for scales smaller than the critical scale Q_c defined by $\det \mathcal{M}^2(Q_c) = 0$ or, equivalently, $m_2^2(Q_c) = 0$. The potential whose minimum, for $Q^2 < Q_c^2$, is given by

$$V_{\text{min}} = - \frac{m_2^4(Q)}{\hat{g}^2} \quad (4)$$

gets deeper as the energy scale Q^2 decreases and the vacuum expectation values are such that $\langle H_1 \rangle = 0$, $\langle H_2 \rangle \neq 0$ (see Figs. 1 and 2). In the reference papers ¹⁰⁾, the physical minimum chosen is the one corresponding to the scale Q_1 defined by the condition

$$\langle H_2 \rangle \equiv \frac{v}{\sqrt{2}}, \quad v \approx 250 \text{ GeV}$$

*) H_2 is the Higgs which gives mass to the top quark.

and the stability condition $m_1^2 + m_2^2 > 0$ is satisfied (see Figs. 1 and 2). However, this tree-level minimum is not the deepest one since deeper minima are generated at scales lower than Q_1 , as we mentioned before, unless supersymmetry is badly broken at scales close to Q_1 . However, we consider this possibility as aesthetically unappealing since the weak scale is introduced by brute force. Disregarding this case, the smallest value of the potential is obtained at the scale Q_0 defined by

$$m_1^2(Q_0) + m_2^2(Q_0) = 0$$

In fact, at this scale, the potential is flat and attains its minimum value along the direction

$$H_2^2 - H_1^2 = - \frac{m_2^2(Q_0)}{\hat{g}^2} \quad (5)$$

(see Figs. 1 and 2). The radiative corrections along this flat direction are important and lift the vacuum degeneracy. Actually this is the line followed in Ref. 11). In the aforementioned mechanisms the condition $m_2^2 < 0$ and hence the symmetry breaking is driven by a large top Yukawa coupling h_t so that they predict a lower bound to the top quark which is $m_t \gtrsim 60$ GeV. Now the following question arises : is it possible to generate a radiative breaking of $SU(2)_L \times U(1)_Y$ with a small top quark mass ? Up to date there are two possible affirmative answers to this question.

- a) Yes, if there exists a fourth generation. This possibility has cosmological troubles (four neutrinos) and perturbative GUT problems (m_b/m_t , large unification coupling constant,...) so that we tend to disregard it.
- b) Yes, if we admit to live in a false vacuum ¹³⁾.

It is clear that none of these solutions is satisfactory. Here we present a new mechanism for radiatively breaking the $SU(2)_L \times U(1)_Y$ in the true vacuum, with only three generations and with small top quark mass. We shall generalize the $SU(2)_L \times U(1)_Y$ radiative breaking for the case $m_3^2 \neq 0$. The choice is not without physical motivation and is dictated to us by the requirement of a correct cosmological behaviour in the low energy phenomenology. The main point is that the photino, $\tilde{\gamma}$, must be the lightest supersymmetric particle. In fact, because $\tilde{\gamma}$ is a Majorana spinor, we have a P-wave suppression in its decay and $m_{\tilde{\gamma}} > O(\text{few GeV})$ ¹⁴⁾. On the other hand, since a neutral Higgsino would have the same suppression factor and a small Yukawa coupling, we deduce ¹⁴⁾ that its mass should be larger than $m_{\tilde{\gamma}}$. However, in models proposed in Refs. 10) and 11),

there is a light Majorana fermion, the axino or supersymmetric partner of the axion, with a mass $m_{\tilde{A}} = (m_3^2/m_{3/2}) \approx 0$ which may violate the cosmological bound. The way out is to have m_3 finite and bounded from below by $m_3^2/m_{3/2} \gtrsim O(\text{few GeV})$. Also another advantage of having a non-zero m_3^2 is the absence of an axion.

In the following, we shall study the radiative breaking mechanism which occurs for non-zero values of m_3^2 . In this case, contrary to the $m_3^2 = 0$ models, the breaking condition given by Eq. (3) does not imply, in general, $m_2^2 < 0$, and the critical scale Q_c is given by $\det \mathcal{M}^2(Q_c) = 0$ or, equivalently, $m_1^2(Q_c)m_2^2(Q_c) = m_3^4(Q_c)$. At scales $Q < Q_c$ the potential has a minimum at

$$\langle H_1 \rangle = \cot \theta(Q) \langle H_2 \rangle \quad (6)$$

and its value at that minimum is given by

$$V_{\min}(Q) = -\frac{1}{\hat{g}^2} \left[\frac{m_1^2 - m_2^2}{2} - \frac{m_1^2 + m_2^2}{2} |\cos 2\theta| \right]^2 \Big|_Q \quad (7a)$$

where

$$\cos^2 2\theta(Q) = \frac{(m_1^2 + m_2^2 + 2m_3^2)(m_1^2 + m_2^2 - 2m_3^2)}{(m_1^2 + m_2^2)^2} \Big|_Q \quad (7b)$$

(see Figs. 3 and 4). As in the $m_3^2 = 0$ case the minimum gets deeper and deeper as the scale Q^2 decreases, attaining its smallest value when $\cos 2\theta = 0$ or, equivalently, when

$$(m_1^2 + m_2^2 - 2m_3^2) \Big|_{Q_0} = 0 \quad (8)$$

The condition given by Eq. (8) defines the scale Q_0 and the minimum of the potential at this scale is

$$V_{\min}(Q_0) = -\frac{1}{\hat{g}^2} \left(\frac{m_1^2(Q_0) - m_2^2(Q_0)}{2} \right)^2 \quad (9)$$

This value is reached when the fields H_1 and H_2 become large and asymptotically equal (see Figs. 3 and 4). This happens because the direction $H_1 = H_2$ is not protected by the D term of the tree-level potential (1) at the scale Q_0 . Also the mass parameter of the field $\phi \equiv (H_1 + H_2)/\sqrt{2}$ is identically equal to zero at Q_0 . Therefore, the radiative corrections along the direction of the field ϕ must be

taken into account and the true vacuum comes from the minimization of the effective potential.

In the leading logarithm approximation, the modification of the tree-level potential along the unprotected ϕ direction is given by

$$\Delta V(\phi) = \frac{1}{64\pi^2} \left[\sum_J (-)^{2J} (2J+1) M_J^4(\phi^2) \right] \log \frac{|\phi|^2}{Q_0^2} \quad (10)$$

The correction along the tree-level protected direction $\chi = (H_1 - H_2)/\sqrt{2}$ is negligible and will not be considered in what follows. The resulting effective potential in terms of ϕ and χ fields turns out to be

$$\begin{aligned} V(\phi, \chi) = & \frac{1}{2} (m_1^2 + m_2^2 - 2m_3^2) |\phi|^2 + \frac{1}{2} (m_1^2 + m_2^2 + 2m_3^2) |\chi|^2 \\ & + \frac{1}{2} (m_1^2 - m_2^2) (\phi \chi^\dagger + \phi^\dagger \chi) + \frac{\hat{g}^2}{2} (\phi \chi^\dagger + \phi^\dagger \chi)^2 + \frac{m^2}{2} |\phi|^2 \log \frac{|\phi|^2}{Q_0^2} \end{aligned} \quad (11)$$

where m^2 is given by

$$m^2 = \frac{d}{d \log Q^2} (m_1^2(Q) + m_2^2(Q) - 2m_3^2(Q)) \Big|_{Q_0} \quad (12)$$

[all the mass parameters in (11) are taken at Q_0]. The fact that $\Delta V(\phi)$ does not contain any $|\phi|^4$ terms is a consequence of the non-renormalization of the D terms because of their non-Abelian structure at energy scales larger than m_X [i.e., the generator of $U(1)_Y$ is traceless]. On the other hand, the specific form of the mass parameter m , Eq. (12), can be traced from the fact that the effective potential satisfies, as a function of ϕ , the same renormalization group equations as the tree-level running mass parameters, as functions of Q .

The minimum of the effective potential (11) is at *)

$$\begin{aligned} \langle \chi \rangle &= -\frac{1}{4} \frac{(m_1^2 - m_2^2) \langle \phi \rangle}{m_3^2 + \hat{g}^2 \langle \phi \rangle^2} \\ \langle \phi \rangle^2 &= \frac{Q_0^2}{e} \exp \left\{ -\frac{(m_1^2 - m_2^2)}{m^2} \frac{\langle \chi \rangle}{\langle \phi \rangle} - 4 \hat{g}^2 \frac{\langle \chi \rangle^2}{m^2} \right\} \end{aligned} \quad (13)$$

*) Condition (8) is assumed and without loss of generality $\langle \chi \rangle$ and $\langle \phi \rangle$ have been taken real.

The fact that $\frac{1}{2}(m_1^2 - m_2^2)$ is smaller than m_{slepton}^2 , as we shall see later, provides the bound (see footnote on p. 8).

$$\left| \frac{\langle \chi \rangle}{\langle \phi \rangle} \right| < \left(\frac{m_{\text{slepton}}}{M_Z} \right)^2 \quad (14)$$

which, in the case of small slepton masses $m_{\text{sl}} = O(20)$ GeV, gives a deviation from the diagonal direction $\langle \chi \rangle = 0$ of the order 5%. Moreover, in the case of small Yukawa couplings, which we will be interested in later, $m_1^2 - m_2^2 = O(m_t^2 m_{\text{slepton}}^2 / M_W^2)$ and the deviation would be much less than 0(1)%. Then the minimum (13) can be written as

$$\begin{aligned} \langle \phi \rangle &\equiv \frac{v}{\sqrt{2}} \simeq \frac{Q_0}{\sqrt{e}} \\ \langle H_1 \rangle &\simeq \langle H_2 \rangle \simeq \langle \phi \rangle / \sqrt{2} \end{aligned} \quad (15)$$

Equation (15) relates the vacuum expectation values $\langle H_1 \rangle, \langle H_2 \rangle$ with the dimensional transmutation scale Q_0 unambiguously. Because of Eq. (15), the contribution of the D terms to the physical mass spectrum of sleptons and squarks is less than 2%. In this way we shall adopt in the following the approximation given by Eq. (15) and neglect the contribution of D terms to the physical mass spectrum. Let us also notice that putting the field χ in its vacuum expectation value, $\langle \chi \rangle \simeq 0$, and solving the renormalization group equation valid for the effective potential, this takes the following compact form :

$$V(\phi) = \frac{1}{2} (m_1^2 (|\phi|^2) + m_2^2 (|\phi|^2) - 2m_3^2 (|\phi|^2)) |\phi|^2 \quad (16)$$

In order to derive the initial conditions at the unification scale of mass parameters and couplings, we define the low energy superpotential as

$$g \equiv h_t U^c Q H_2 + m_4 H_1 H_2 \quad (17)$$

where U^c stands for the chiral superfield which contains the right-handed top quark, Q the one containing the left-handed top-bottom doublet and H_1, H_2 the superfields containing the usual Higgs scalars ^{*)}. We are neglecting all

^{*)} H_1, H_2 now denote the complete Higgs doublets and not just their neutral components as in Eq. (1).

the other Yukawa couplings as compared to the top-quark Yukawa coupling h_t^* . The resulting scalar potential, including the soft breaking terms coming from supergravity and relative to the superfields U^c , Q and H_1 , is given, at scales $\approx M_X$, by

$$\begin{aligned} V = & m_{3/2}^2 (|Q|^2 + |U^c|^2) + (m_{3/2}^2 + m_4^2) (|H_1|^2 + |H_2|^2) \\ & + (A-1) m_{3/2} m_4 (H_1 H_2 + h.c.) + h_t A m_{3/2} (U^c Q H_2 + h.c.) \\ & + m_4 (H_1^\dagger U^c Q + h.c.) + \text{"D terms"} \end{aligned} \quad (18)$$

where $m_{3/2}$ is the gravitino mass and we are using the same symbol to denote a superfield and its scalar component. One then obtains the following initial conditions for m_1^2 , m_2^2 , m_3^2 , m_Q^2 and m_U^2 :

$$\begin{aligned} m_1^2(M_X) &= m_2^2(M_X) = m_{3/2}^2 + m_4^2(M_X) \\ m_Q^2(M_X) &= m_U^2(M_X) = m_{3/2}^2 \\ m_3^2(M_X) &= (A(M_X) - 1) m_{3/2} m_4(M_X) \equiv B(M_X) m_{3/2} m_4(M_X) \end{aligned} \quad (19)$$

Taking into account the renormalization group equations for m_1^2 , m_2^2 , m_3^2 , A , B and h_t as well as for the gauge couplings and gaugino masses, one obtains the following equations at the Q_0 scale

$$\begin{aligned} m_1^2 &= m_{sl}^2 + m_4^2 \\ m_2^2 &= m_{sl}^2 + m_4^2 - 3r^2 (m_2^2 - m_4^2 + m_U^2 + m_Q^2 + m_{3/2}^2 A^2) \\ m_3^2 &= B m_{3/2} m_4 \\ m_Q^2 &= m_{3/2}^2 (1 + C_Q \xi^2) - \frac{1}{3} (m_1^2 - m_2^2) \\ m_U^2 &= m_{3/2}^2 (1 + C_U \xi^2) - \frac{2}{3} (m_1^2 - m_2^2) \\ B &= A(M_X) - 1 + C_B \xi - 3r^2 A \\ A &= (A(M_X) + C_A \xi) / (1 + 6r^2) \end{aligned} \quad (20)$$

*) This approximation is obviously valid in our case where $\langle H_1 \rangle \approx \langle H_2 \rangle$.

where ξ is the ratio of the gaugino to the gravitino mass at M_X and m_{sl} stands for the mass of the left-handed sleptons

$$m_{sl}^2 = m_{3/2}^2 (1 + C_L \xi^2) \quad (21)$$

C_A, C_B, C_U, C_Q and C_L are model-dependent parameters given as integrals of powers of gauge coupling constants and r^2 is the integral

$$r^2 \equiv \int_0^{\log M_X^2/M_W^2} \left(\frac{h_t}{4\pi} \right)^2 d \log \frac{Q^2}{M_W^2} \quad (22)$$

Actually, Eqs. (20) are approximate analytic solutions ⁷⁾ to the renormalization group equations valid for a large class of models ^{6),7),10),11)}. The top quark mass is given by ⁷⁾

$$\left(\frac{m_t}{M_W} \right)^2 \simeq 19 \frac{r^2}{1+6r^2} \quad (23)$$

Using the condition $(m_1^2 + m_2^2 - 2m_3^2)|_{Q_0} = 0$, Eq. (8) and Eqs. (20), we obtain the algebraic equation for r^2

$$r^2 = \frac{2}{3} \frac{m_{sl}^2 - m_4(Bm_{3/2} - m_4)}{m_{3/2}^2 [-1 + (C_U + C_Q - 3C_L)\xi^2 + A^2] + 4m_4(Bm_{3/2} - m_4)} \Big|_{M_W} \quad (24)$$

The minimum value of r^2 , i.e., the minimum value of m_t , is reached when $m_4 = m_3 B/2$ ^{*)}. We shall assume this condition in the following since we are really interested for small values of m_t . The stability condition at m_X given by Eq. (2) restricts the parameters to satisfy the bound

$$\left[m_{3/2}^2 - m_{3/2} m_4 (A-1) + m_4^2 \right] \Big|_{M_X} > 0 \quad (25)$$

Also, in order to avoid minima which break colour and/or electromagnetism to all energy scales from M_X down to M_W the following bounds should be satisfied ⁷⁾

^{*)} This actually provides the initial value of m_4 at M_X once m_3 and $A(M_X)$ are given. Notice that for $m_3^2 \equiv m_4 m_3 B < m_4^2$, one obtains through Eqs. (23) and (24) values for m_t larger than those of Ref. 11). We therefore assume that $m_3^2 > m_4^2$ which leads to the bound of Eq. (14).

$$|A(M_X)|^2 < 3 \left(3 + \frac{m_4^2}{m_{3/2}^2} \right) \Big|_{M_X} \quad (26a)$$

$$|A(M_W)|^2 < 3 \left(\frac{m_Q^2 + m_U^2 + m_2^2}{m_{3/2}^2} \right) \Big|_{M_W} \quad (26b)$$

Actually, these bounds are more restrictive than those given in Refs. 10) and 11). Notice that (25) and (26a) are automatically satisfied in the range $-1 \leq A(M_X) \leq 3$. Owing to these bounds " r^2 " as given by Eq. (24) cannot become arbitrarily small in the $m_4 = 0$ ($m_3^2 = 0$) case. This, in conjunction with Eq. (23), yields in the $m_3^2 = 0$ models the well-known lower bound of 60 GeV for m_t . It is obvious that if one relaxes the bounds (26) m_t^2 can become small as $|A|$ becomes large at the cost of having metastable $SU(2)_L \times U(1)_Y$ breaking minima 11), 13).

In the following, we shall present our results for two particular models :

- a) the minimal model (MIM) with three families and two Higgs doublets, and
- b) the cosmologically acceptable supergravity model (CAM) 6).

In the CAM, two natural intermediate scales appear $\Lambda_1 \approx 10^{4-6}$ GeV and $\Lambda_2 \approx 10^{10-12}$ GeV. The particle content is : four colour triplets of mass $\sim \Lambda_2$; one colour octet and one $SU(2)$ triplet of mass $\sim \Lambda_1$; four light ($\sim M_W$) singlets originating from {10} representations of $SU(5)$, four light ($\sim M_W$) doublets and three generations of matter field.

For the MIM, the constants $C_{A,B,Q,U,L}$ appearing in Eqs. (20)-(24) are given by 7)

$$\begin{aligned} C_A &= 3.30, & C_B &= .55 \\ C_Q &= 5.29, & C_U &= 4.89, & C_L &= .50 \end{aligned} \quad (27a)$$

while for the CAM the corresponding constants take the following values 7)

$$\begin{aligned} C_A &= 2.38, & C_B &= .56 \\ C_Q &= 1.82, & C_U &= 1.53, & C_L &= .36 \end{aligned} \quad (27b)$$

In both models, the smallest value of m_t is obtained when $A(M_X)$ is between 2.5 and 3 provided that $\xi \equiv (M_{\text{gaugino}}/m_{3/2}) M_X$ takes reasonable values $\lesssim 0(10)$. In the Table we present a typical physical spectrum of the two models for $m_t = 25$ GeV, $m_t = 35$ GeV and $m_t = 50$ GeV, respectively. The gravitino mass in each case has been chosen so as to reach the lowest spectrum compatible with present-day experimental bounds.

A few comments are in order, which concern the spectra of the two models.

For the same value of the parameter ξ :

- i) the squarks in the MIM are heavier than in the CAM ;
- ii) the slepton spectrum is almost identical in both models ;
- iii) the gluino in the MIM is more than twice as heavy as in the CAM ; actually

$$\frac{M_{\text{gluino}}(\text{MIM})}{M_{\text{gluino}}(\text{CAM})} = \frac{\alpha_{G47}(\text{CAM})}{\alpha_{G47}(\text{MIM})} \approx 2-3.$$

- iv) for light top quark and slepton masses the photino mass in the MIM may just saturate the lower cosmological bound unlike the cases in the CAM where this bound is obviously satisfied ;
- v) the supersymmetric partner of the would-be axion (axino) is always heavier than the photino, in both models, in agreement with cosmological bounds.
- vi) a common characteristic of these models is the existence of a "light" neutral Higgs boson ($m_{H_0} \lesssim 10$ GeV) which may easily show up in heavy Onium [including $Y(b\bar{b})$] decays like $0 \rightarrow H^0 + \gamma$.

In conclusion, we have shown that it is possible to keep the very interesting and appealing idea of radiative breaking of $SU(2)_L \times U(1)_Y$ by supergravity effects without necessarily invoking large top quark masses or unstable minima. At the same time, we provide a realistic and experimentally accessible low energy particle spectrum.

PARTICLE SPECTRUM

		CAM	MIM	CAM	MIM	CAM	MIM
$A(M_X)$		3	3	2.8	2.8	2.0	1.6
$m_{\frac{3}{2}}$		15	15	15	15	15	15
ξ		2.8	2.2	3.2	3.1	3.5	1.9
$m_u(M_X)$		15	17	16	18	11	7
top		25	25	35	35	50	50
All families	(sleptons) _L	29	27	32	36	35	25
	(sleptons) _R	21	20	22	23	23	19
1st and 2nd families	(squark) _L	58	77	66	108	72	67
	(squark) _R }	54	74	61	104	67	65
		54	74	60	103	66	65
3rd family	(sbottom) _L	58	76	64	106	68	66
	(sbottom) _R	54	74	60	104	66	65
	(stop) _L	81	96	95	132	112	106
	(stop) _R	26	54	23	78	21	37
Charged Higgses		96	93	95	94	88	83
Neutral Higgses }		106	104	105	105	100	95
		3	3	4.4	5.3	6	5
"Axion"		51	46	49	48	35	19
Gluinos		42	84	47	118	52	72
Photino		11	4.6	9.4	7.3	4	3
HW, WH-inos }		89	87	87	94	84	90
		78	82	79	79	82	75
HZ, HZ-inos }		99	108	99	116	101	106
		92	85	91	80	88	83
Axino		26	23	24	24	17	9

Physical mass spectrum of the cosmologically acceptable model (CAM) and the minimal model (MIM) corresponding to the same gravitino mass $m_{\frac{3}{2}} \approx 15$ GeV for top quark masses equal to 25 GeV, 35 GeV and 50 GeV, respectively. ξ denotes the ratio of the gaugino to the gravitino mass at M_X . All masses are in GeV units. The light neutral Higgs gets its mass via radiative corrections.

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FIGURE CAPTIONS

Fig. 1 Evolution of $\langle H_2 \rangle$, $V_{\min}$ and $m_{1,2}^2$ as functions of the scale Q^2 for $m_3^2 = 0$. The critical point Q_C^2 is shown as well as the scale Q_0^2 defined by $m_1^2(Q_0) + m_2^2(Q_0) = 0$. $V_{\min}(Q)$ obtains its minimum value at Q_0 .

Fig. 2 $V_{\min}$ as a function of the vacuum expectation values $\langle H_1 \rangle$ and $\langle H_2 \rangle$ for different scales Q^2 when $m_3^2 = 0$. For $Q^2 > Q_C^2$, $V_{\min} = 0$. The deepest $V_{\min}$ is obtained along the flat direction $\langle H_2 \rangle^2 - \langle H_1 \rangle^2 = -m_2^2(Q_0)/\hat{g}^2$.

Fig. 3 Evolution of $\langle H_1 \rangle$, $\langle H_2 \rangle$, $V_{\min}$, $m_1^2 + m_2^2$ and m_3^2 as functions of the scale Q^2 when $m_3^2 \neq 0$. The critical point Q_C satisfies $m_1^2(Q_C)m_2^2(Q_C) = m_3^4(Q_C)$ and the scale Q_0 is defined by $m_1^2(Q_0) + m_2^2(Q_0) = 2m_3^2(Q_0)$. For $Q^2 > Q_C^2$, $\langle H_{1,2} \rangle = 0$. For scales $Q_0^2 < Q^2 < Q_C^2$ $\langle H_1 \rangle = \langle H_2 \rangle \cot\theta(Q)$.

Fig. 4 $V_{\min}$ as a function of the vacuum expectation values $\langle H_1 \rangle$ and $\langle H_2 \rangle$ for different scales Q^2 in the case $m_3^2 \neq 0$. For $Q^2 > Q_C^2$, $V_{\min} = 0$. The deepest $V_{\min}$ occurs for $\langle H_1 \rangle$, $\langle H_2 \rangle$ along the curve $\langle H_1 \rangle = \langle H_2 \rangle \cot\theta(Q)$. The angle $\theta(Q)$ is defined in the text [see Eq. (7b)].

FIG. 1

FIG. 2

FIG. 3

FIG. 4