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One-loop analysis of the electroweak breaking in supersymmetric models and the fine-tuning problem

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We examine the electroweak breaking mechanism in the minimal supersymmetric standard model (MSSM) using the *complete* one-loop effective potential V_1. First, we study what is the region of the whole MSSM parameter space (i.e. $M_{1/2}, m_0, \mu, ...$) that leads to a successful $SU(2) \times U(1)$ breaking with an acceptable top-quark mass. In doing this it is observed that all the one-loop corrections to V_1 (even the apparently small ones) must be taken into account in order to obtain reliable results. We find that the allowed region of parameters is considerably enhanced with respect to former "improved" tree-level results. Next, we study the fine-tuning problem associated with the high sensitivity of M_Z to h_t (the top Yukawa coupling). Again, we find that this fine-tuning, once the one-loop effects are considered, is appreciably smaller than in previous tree-level calculations. Finally, we explore the ambiguities and limitations of the ordinary criterion to estimate the degree of fine-tuning. As a result of all this, the upper bounds on the MSSM parameters, and hence on the supersymmetric masses, are substantially raised, thus increasing the consistency between supersymmetry and observation.

1. Introduction

Precision measurements at LEP give a strong support [1] to the expectations of supersymmetric (SUSY) [2] grand unification [3]. Namely, the two-loop calculation indicates that the gauge coupling constants of the standard model seem to be unified [#1] at $M_X \sim 10^{16}$ GeV with a value $\alpha_X \sim 1/26$, provided the average mass of the new supersymmetric states lies in the range [100 GeV, 10 TeV].

This calculation has been refined in a recent paper by Ross and Roberts [5] in which the various supersymmetric thresholds were appropriately taken into account. This was done in the context of the minimal supersymmetric standard model (MSSM), which is characterized by the Lagrangian

$$\mathcal{L} = \mathcal{L}_{\text{SUSY}} + \mathcal{L}_{\text{soft}} . \tag{1}$$

Here $\mathcal{L}_{\text{SUSY}}$ is the supersymmetric Lagrangian derived from the observable superpotential W_{obs}, which includes the usual Yukawa terms W_Y and a mass coupling $\mu H_1 H_2$ between the two Higgs doublets H_1, H_2. At the unification scale M_X, $\mathcal{L}_{\text{soft}}$ is given by

$$\begin{aligned}\mathcal{L}_{\text{soft}} &= -m_0^2 \sum_\alpha |\phi_\alpha|^2 - \tfrac{1}{2} M_{1/2} \sum_{a=1}^{3} \bar{\lambda}_a \lambda_a \\ &\quad - (A m_0 W_Y + B m_0 \mu H_1 H_2 \; + \; \text{h.c.}) ,\end{aligned} \tag{2}$$

where m_0 and $M_{1/2}$ are the (common) supersymmetry soft breaking masses (at M_X) for all the scalars ϕ_α and gauginos λ_a of the theory, and A and B parametrize the (common) couplings of the trilinear and bilinear scalar terms. In this framework the physical spectrum of supersymmetric masses depends on the particular choice of the MSSM parameters

$$m_0, \; M_{1/2}, \; \mu, \; A, \; B, \; h_t , \tag{3}$$

[#1] This unification does not necessarily require a GUT. In particular, in superstring theories all the gauge couplings are essentially the same at tree level [4], even in the absence of a grand unification group. This also avoids unwanted consequences of GUT theories.

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where h_t is the top Yukawa coupling [#2]. Therefore, the requirement of gauge unification constrains their ranges of variation.

These parameters are also responsible for the form of the Higgs scalar potential and thus for the electroweak breaking process [6]. Requiring the electroweak scale (i.e. M_Z) to be the correct one, together with the present bounds on m_t, Ross and Roberts further restricted the allowed space of these parameters. Finally, the authors imposed the absence of fine-tuning in the value of h_t (the parameter to which M_Z is most sensitive [#3]) for a successful electroweak breaking, by demanding [7] $c \lesssim 10$ in the equation

$$\frac{\delta M_Z^2}{M_Z^2} = c\frac{\delta h_t^2}{h_t^2}, \tag{4}$$

where the value of c depends on the values of all the independent parameters shown in (3) (which also determine the supersymmetric masses). As a consequence, they found $m_0, \mu, M_{1/2} \lesssim 200$ GeV (leading to typical supersymmetric masses $\lesssim 500$ GeV). In fact, this turns out to be the strongest constraint on the supersymmetric mass scale, stronger than the requirement of gauge unification.

The analysis of ref. [5] of the electroweak breaking process and the corresponding h_t-fine-tuning problem was performed by using the renormalization-improved tree-level potential $V_0(Q)$, i.e. the tree-level potential in terms of the renormalized parameters at the scale Q. However, as was shown in ref. [8], the effect of the one-loop contributions is expected to be important [#4]. Consequently, the analysis should be redone using the whole one-loop effective potential. This is the main goal of this paper.

In section 2 we study what is the region of the whole MSSM parameter space (formula (3)) leading to a correct $SU(2) \times U(1)$ breaking (this means a correct value for M_Z and m_t without colour and electric charge breakdown). The comparison with the results of the "renormalization-improved" tree-level potential V_0 [5] shows that the one-loop corrections enhance (and also displace) this allowed region. As a by-product, we show that the (very common) approximation of considering only the top and stop contribution (disregarding the $\tilde{t}_L - \tilde{t}_R$ mixing) to the one-loop effective potential is not reliable for analysing the electroweak breaking mechanism. In section 3 we analyse the above-mentioned fine-tuning problem, showing that, once the one-loop contributions are taken into account, it becomes considerably softened. In addition to this, we study the limitations and ambiguities of the ordinary criterion (4) to estimate the fine-tuning problem. Although in the MSSM it turns out to be a sensible criterion (which is not a general fact), it should be considered as a rather qualitative one; the upper bound on c should thus be conservatively relaxed, at least up to $c \lesssim 20$. As a consequence of all this, the upper bounds on the MSSM parameters and on the supersymmetric masses are pushed up from the "renormalization-improved" tree-level results. This is relevant, of course, for the expectations of experimental detection of SUSY. Finally, we present our conclusions in section 4.

2. Radiative electroweak breaking

In the MSSM the part of the tree-level potential along the neutral components of the Higgs fields at a scale Q is given by

$$V_0(Q) = \tfrac{1}{8}(g^2 + g'^2)\left(|H_1|^2 - |H_2|^2\right)^2 + m_1^2|H_1|^2 + m_2^2|H_2|^2 - m_3^2(H_1H_2 + \text{h.c.}), \tag{5}$$

where

$$m_i^2 = m_{H_i}^2 + \mu^2, \quad m_3^2 = m_0\mu B, \tag{6}$$

with

$$m_{H_i}^2(M_X) = m_0^2. \tag{7}$$

In the usual calculations with just the tree-level potential $V_0(Q)$ (as in ref. [5]), this was minimized at the M_Z (or M_W) scale.

The one-loop effective potential is given by [10]

$$V_1(Q) = V_0 + \Delta V_1, \tag{8}$$

where

[#2] These are the parameters, together with the gauge couplings, that enter in the renormalization group equations for the masses. The influence of the bottom and tau Yukawa couplings is negligible in most of the cases.

[#3] The sensitivity of M_Z to other independent parameters has been analysed in ref. [7].

[#4] For recent work on this subject see ref. [9].

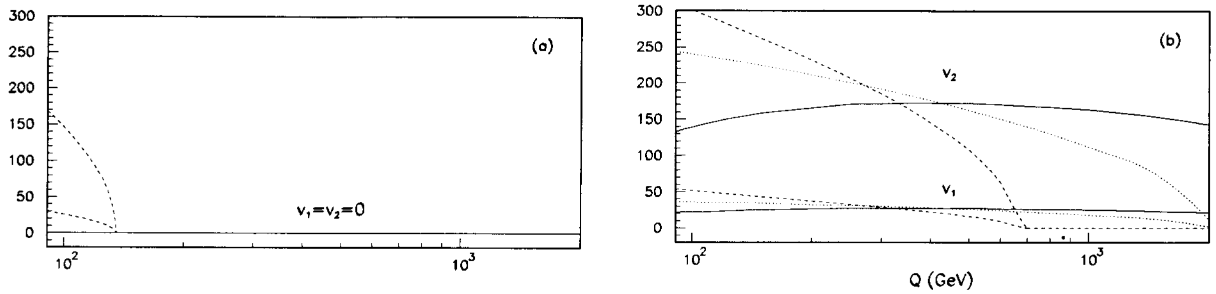

Fig. 1. $v_1 \equiv \langle H_1 \rangle$, $v_2 \equiv \langle H_2 \rangle$ versus the Q scale between M_Z and 2 TeV (in GeV) for the cases labelled (a) and (b) in eq. (10). Solid lines: complete one-loop results; dashed lines: "improved" tree-level results; dotted lines: one-loop results in the top–stop approximation.

$$\Delta V_1(Q) = \frac{1}{64\pi^2}\mathrm{Str}\left[\mathcal{M}^4\left(\log\frac{\mathcal{M}^2}{Q^2} - \tfrac{3}{2}\right)\right], \tag{9}$$

depends on H_1, H_2 through the tree-level squared-mass matrix $\mathcal{M}^2$. In the expressions (5), (6), (8), (9) all the parameters are understood to be running parameters evaluated at the scale Q. They can be computed by solving the standard renormalization group equations (RGEs), whose form is well known [2], and taking into account all the supersymmetric thresholds. The supertrace of eq. (9) runs over *all* the states of the theory. This, in particular, amounts to determine the eigenvalues of the mass-mixing matrices of stops, charginos and neutralinos. Incidentally, a simplification broadly used in the literature is to consider just the top (t) and stop ($\tilde{t}$) contributions to (9), disregarding the $\tilde{t}_L - \tilde{t}_R$ mixing. This can be a good approximation for certain purposes (see e.g. ref. [11]), but *not*, as will be shown shortly, when one is interested in studying the SU(2) × U(1) breaking. To be on the safe side, the whole spectrum contribution must be considered in eq. (9).

In order to exhibit the implications of considering the whole one-loop potential V_1 versus V_0, we have shown two examples (a) and (b) in fig. 1. They are specified by the following initial values of the independent parameters

$$\begin{aligned}
&\text{(a)} \quad m_0 = \mu = 120\ \mathrm{GeV}, \quad M_{1/2} = 230\ \mathrm{GeV}, \\
&\qquad A = B = 0, \quad h_t = 0.207, \\
&\text{(b)} \quad m_0 = \mu = 100\ \mathrm{GeV}, \quad M_{1/2} = 180\ \mathrm{GeV}, \\
&\qquad A = B = 0, \quad h_t = 0.250.
\end{aligned} \tag{10}$$

The case (a) corresponds to one of the two models explicitly expounded in ref. [5] (where it was called X). Although in the V_0 approximation this model works correctly, once the one-loop contributions are considered, we see that it does not even lead to electroweak breaking (the same happens with the model that was called Z). In the example (b), both V_0 and V_1 yield electroweak breaking, but for completely different values of $v_1 \equiv \langle H_1 \rangle$ and $v_2 \equiv \langle H_2 \rangle$. In this case, V_1 predicts electroweak breaking at the right scale, while V_0 does not. The above-mentioned approximation of considering just the top and stop contribution to ΔV_1, which is also represented in the figure, works better than V_0, but not enough to produce acceptable results. Moreover, it is clear from the figure that only the *whole* one-loop contribution really helps to stabilize the values of v_1, v_2 versus variations of Q (they are essentially constant up to $\mathrm{O}(\hbar^2)$ corrections). In fact, they should evolve only via the (very small) wave function renormalization effects, given by

$$\begin{aligned}
\frac{\partial \log v_1}{\partial \log Q} &= \frac{1}{64\pi^2}(3g_2^2 + g'^2), \\
\frac{\partial \log v_2}{\partial \log Q} &= \frac{1}{64\pi^2}(3g_2^2 + g'^2 - 12h_t^2).
\end{aligned} \tag{11}$$

There is a scale, which in ref. [8] was called $\widehat{Q}$, at which the results from V_0 and V_1 approximately coincide. At this scale the one-loop contributions are quite small, in particular the logarithmic factors, so $\widehat{Q}$ represents a certain average of all the masses. In the region around $\widehat{Q}$ one expects, because of the smallness of the logarithms, that the evaluation of one-loop effects is more reliable (see also ref. [12]).

In the example depicted in fig. 1b this consideration is not very relevant, for v_1 and v_2 are essentially

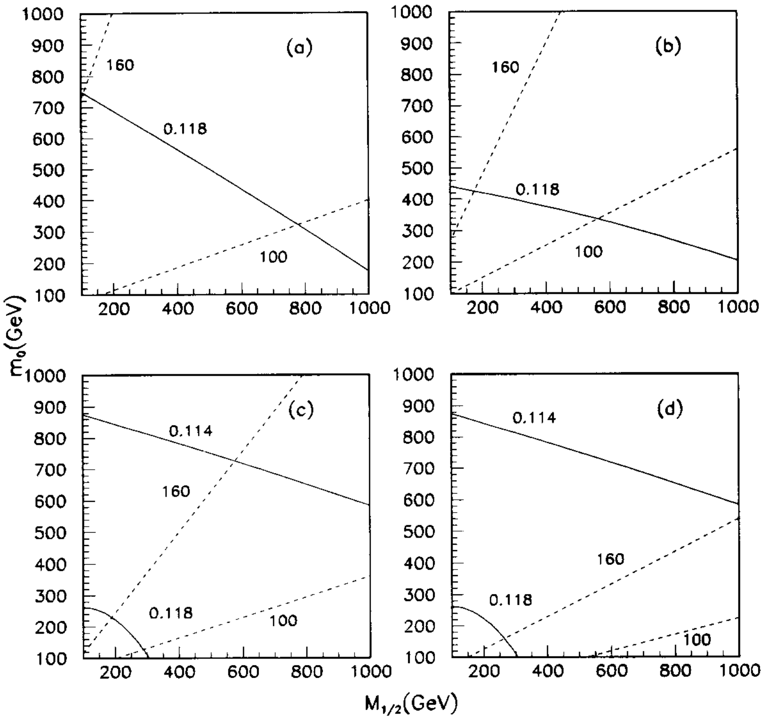

Fig. 2. Allowed values for the $M_{1/2}$, m_0 parameters (in GeV) for different values of μ_0: $|\mu_0/m_0| = 0.2, 0.4, 1, 3$ in (a), (b), (c), (d) respectively, and $A = B = 0$. The solid lines represent the value of $\alpha_3(M_Z)$ needed to achieved unification, as calculated in ref. [5]. Dotted lines correspond to the extreme values of m_t (evaluated at the M_Z scale): $m_{\text{top}} = 160, 100$ GeV.

constant. However, there are cases where $v_1(Q)$ and $v_2(Q)$ do not show such a remarkable stability. This happens when the averaged supersymmetric mass is much larger than M_Z, since this leads to the appearance of large logarithms at $Q = M_Z$ (this fact has been stressed in ref. [12]). However, in the region around $\widehat{Q}$ (i.e. precisely where the calculation is more reliable) $v_1(Q)$ and $v_2(Q)$ are *always* stable. Thus we have used the following criterion: we evaluate v_1 and v_2 at the $\widehat{Q}$ scale and then calculate $v_1(Q)$ and $v_2(Q)$ via eq. (11) at any other scale. This is relevant at the time of calculating physical masses. In particular, M_Z is given by

$$(M_Z^{\text{phys}})^2 \simeq \tfrac{1}{2}\left(g_2^2(Q) + g'^2(Q)\right) \times \left[v_1^2(Q) + v_2^2(Q)\right]\big|_{Q=M_Z^{\text{phys}}} \tag{12}$$

and similar expressions can be written for all the particles of the theory.

Now we are ready to determine how the requirement of correct electroweak breaking puts restrictions on the space of parameters. "Correct electroweak breaking" of course means $M_Z^{\text{phys}} = M_Z^{\text{exp}}$, where M_Z is given by eq. (12). In addition, other physical requirements must be satisfied. Namely, the scalar potential must be bounded from below [2], colour and electric charge must remain unbroken [2], and the top mass must lie within the LEP limits (100 GeV $\lesssim m_{\text{top}} \lesssim$ 160 GeV). Following a presentation similar to that of ref. [5], the results of the analysis for $A = B = 0$ (at M_X) and for various initial values of $|\mu_0/m_0|$ are shown in fig. 2. The value of $\alpha_3(M_Z)$ necessary to achieve unification of the couplings was calculated in ref. [5] at the two-loop order and is also represented in the figure. We have also evaluated the effect of varying the A and B parameters, as is illustrated in fig. 3. The effect of the one-loop contribution is to enhance and displace the region of allowed parameters appreciably. In order

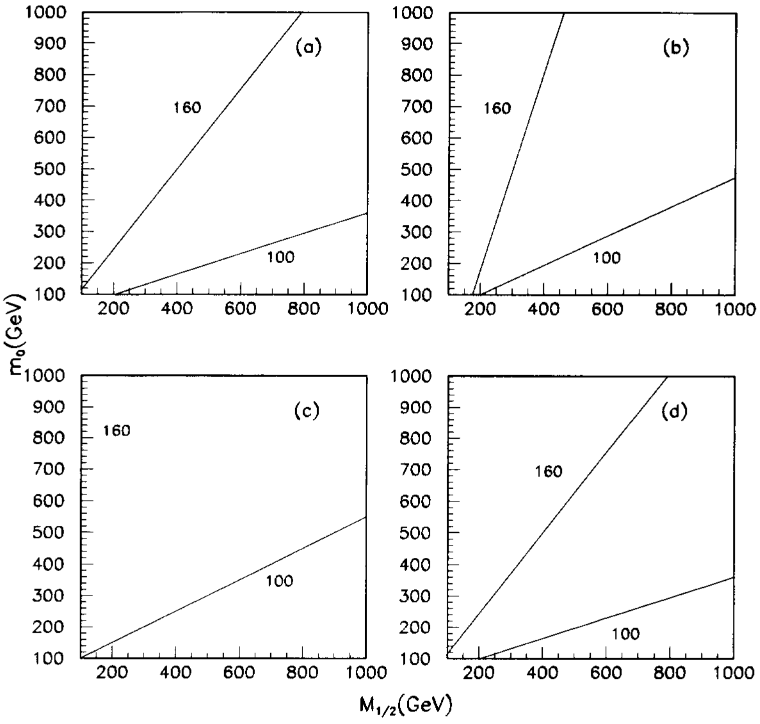

Fig. 3. The same as fig. 2, but for different values of A, B: $A = 0, 0, 1, -1$, $B = 0, 1, 0, 0$ in (a), (b), (c), (d) respectively, and $|\mu_0/m_0| = 1$. In case (c), the $m_t = 160$ GeV line coincides with the $M_{1/2} = 100$ GeV axis.

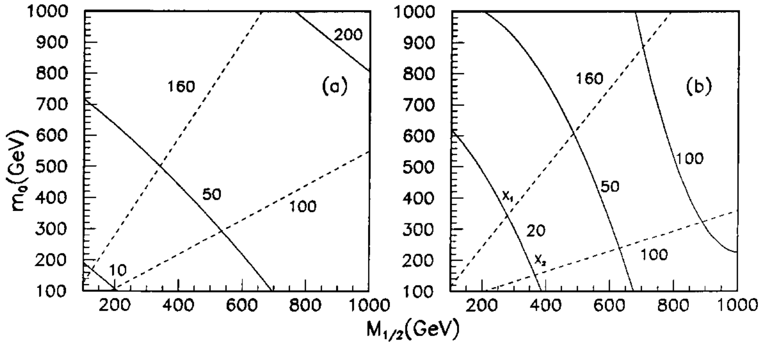

Fig. 4. The case $A = B = 0$, $|\mu_0/m_0| = 0.2, 0.4, 1, 3$ with the "improved" tree-level potential V_0 (a) and the whole one-loop effective potential V_1 (b). Diagonal lines correspond to the extreme values of m_t, as were calculated by Ross et al. in ref. [5]: $m_t = 160, 100$ GeV. Transverse lines indicate constant values of c, defined in eq. (4).

to facilitate the comparison we have reproduced in fig. 4 the V_0 results [5] and the one-loop results for the case of fig. 2c (i.e. $m_0/\mu_0 = 1, A = B = 0$), which is a representative one.

3. The fine-tuning problem

As was pointed out in ref. [5], h_t is the parameter to which the value of M_Z is more sensitive. This sensitivity is conveniently quantified by the c parameter defined in eq. (4). We have represented the values of c for the representative case of fig. 4. A good parametrization of the value of c is

$$c \simeq \frac{1}{M_Z^2}\left[1.08M_{1/2}^2 + 0.19\,(m_0^2 + \mu_0^2)\right] . \tag{13}$$

The strong influence of $M_{1/2}$ on the value of c compared with that of m_0 and μ_0 comes from the fact that scalar masses can be very high, even if they are vanishing at tree level, owing to the gaugino contribution in the RGEs, but not the other way round. The tree level results [5] are also given to facilitate the comparison [#5]. The sensitivity of M_Z to h_t turns out to be substantially smaller with the complete one-loop effective potential than with the V_0 approximation. If, following ref. [5], we now demand $c \lesssim 10$ as the criterion to avoid the fine-tuning in h_t, this selects a region of acceptable SUSY parameters that can easily be read from fig. 4. This region is clearly larger than the corresponding one obtained from V_0. This is a consequence of the lower sensitivity of M_Z to h_t *and* of the larger region of parameters giving a correct value of M_Z (see section 2) when one uses the entire one-loop effective potential V_1. Accordingly, the one-loop contributions tend to make the electroweak breaking process in supersymmetric models less "critical".

We would also like to make some comments on the criterion usually followed to parametrize the fine-tuning problem, i.e. $c \lesssim 10$ in eq. (4). First of all, to some extent this procedure is ambiguously defined, since it depends on our definition of the independent parameters and the physical magnitude to be fitted. For example, if we replace M_Z^2 by M_Z in eq. (4), then the corresponding values of c (represented in fig. 3) are divided by two. Secondly, notice that if for a certain choice of the supersymmetric parameters $(m_0, M_{1/2}, \mu, A, B)$, the value of c turned out to be high for most of the possible values of h_t (or equivalently M_Z), then we would arrive at the bizarre conclusion that *any* value of h_t leads to a fine-tuning [#6]! This is so because the "standard" criterion of eq. (4) measures the *sensitivity* of M_Z to h_t rather than the degree of fine-tuning. In order for eq. (4) to be a sensible quantification of the fine-tuning, it should be required that $c \sim 1$ for most of the h_t values. To check this, we have represented in fig. 5 M_Z versus h_t for a typical example ($m_0 = \mu = M_{1/2} = 500$ GeV, $A = B = 0$). We see that, indeed, for most of the h_t values the sensitivity of M_Z to h_t is small. Hence, the parametrization of the fine-tuning by the value of c in eq. (4) is meaningful. A natural value for M_Z under these conditions would be $M_Z \sim 1$ TeV [#7]. Nevertheless, all this shows that it is dangerous to assume that c is an exact measure of the degree of fine-tuning. It is rather a sensitive, but qualitative one. In fact, a precise evaluation of the degree of fine-tuning would require a knowledge of what the actual independent parameters of the theory are and what the supergravity breaking mechanism is (for an example of this see ref. [13]).

All the previous considerations suggest that the upper limit $c \lesssim 10$ in the measure of the allowed fine-tuning should be conservatively relaxed, at least up to $c \lesssim 20$. We see from fig. 4 that this implies

$$m_0, \mu \lesssim 650 \text{ GeV}, \quad M_{1/2} \lesssim 400 \text{ GeV}. \tag{14}$$

In order to see what the corresponding upper limits on the supersymmetric masses are, we have explicitly given the mass spectrum (including also the small contributions coming from the electroweak breaking) in table 1 for the two "extreme" cases labelled X_1 and X_2 in fig. 4. Note that these two cases are close to the $c = 20$ line and to the upper and lower limits on the top-quark mass. From these extreme examples we see

#5 We reproduce here the values of c for V_0 as given in ref. [5], although our calculation gives slightly different values.

#6 This would happen, for instance, if the hypothetic theoretical relation between M_Z and h_t were $M_Z \sim \exp\{Ch_t\}$ with $|Ch_t| > 10$.

#7 Notice, however, that if we restrict the range of variation of h_t so that 100 GeV $< m_t <$ 160 GeV, then $c > 10$ in the entire "allowed" region of h_t.

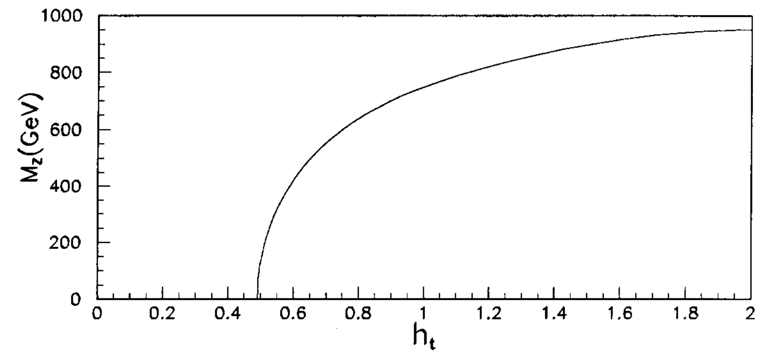

Fig. 5. M_Z versus h_t for $M_{1/2} = m_0 = \mu_0 = 500$ GeV, $A = B = 0$. The region of physical M_Z amounts to a fine-tuning in the value of h_t.

Table 1
Masses of the supersymmetric states for the two solutions (called X_1 and X_2 in fig. 4) with $m_t = 163$, 109 GeV respectively. All the masses are given at the M_Z scale.

Parameters (initial values)		
$M_{1/2}$ (GeV)	300	400
m_0 (GeV)	400	200
μ (GeV)	400	200
h_t	0.618	0.254
A, B	0	0
Masses of gluino, charginos and neutralinos (in GeV)		
$\tilde{g}$	837	1124
$\chi_1^{\pm}$	407	376
$\chi_2^{\pm}$	243	226
λ_1	172	169
λ_2	242	371
λ_3	408	236
λ_4	387	255
Masses of squarks (in GeV)		
$\tilde{u}_L, \tilde{c}_L$; $\tilde{d}_L, \tilde{s}_L$	785; 789	922; 925
$\tilde{u}_R, \tilde{c}_R$	766	888
$\tilde{d}_R, \tilde{s}_R, \tilde{b}_R$	762	885
$\tilde{t}_L, \tilde{b}_L$	827, 698	1055, 881
$\tilde{t}_R$	410	560
Masses of sleptons and higgses (in GeV)		
$\tilde{l}_L, \tilde{l}_R$	476, 431	372, 256
h^0, H^0	91, 547	91, 353
$H^{\pm}$	553	362
A^0	547	353

that, roughly speaking, the bounds on the most relevant supersymmetric particles are

$$\begin{array}{ll} \text{Gluino} & M_{\tilde{g}} \lesssim 1100\ \text{GeV}, \\ \text{Lightest chargino} & M_{\chi^{\pm}} \lesssim 250\ \text{GeV}, \\ \text{Lightest neutralino} & M_{\lambda} \lesssim 200\ \text{GeV}, \\ \text{Squarks} & m_{\tilde{q}} \lesssim 900\ \text{GeV}, \\ \text{Sleptons} & m_{\tilde{l}} \lesssim 450\ \text{GeV}. \end{array} \tag{15}$$

These numbers are substantially higher than those obtained in ref. [5] from V_0, and summarize the three main results obtained in this paper: (i) The region of parameters giving a correct electroweak breaking is larger when one uses the entire one-loop effective potential V_1 than with V_0 (see section 2). (ii) The corresponding sensitivity of M_Z to the value of h_t is smaller. (iii) The highest acceptable value of c (see eq. (4)) must be conservatively relaxed for the reasons explained above. The most important conclusion at this stage is that the supersymmetric spectrum is not necessarily close to the present experimental limits, although the future accelerators (LHC, SSC) should bring it to light. It is also remarkable that the $\tilde{t}_L - \tilde{t}_R$ splitting can be sizeable in many scenarios. Let us finally note that there are considerable radiative corrections to the lightest-Higgs mass coming from the top–stop splitting [12], which have not been included in table 1.

4. Conclusions

We have studied the electroweak breaking mechanism in the minimal supersymmetric standard model (MSSM) using the *complete* one-loop effective potential $V_1 = V_0 + \Delta V_1$ (see eqs. (5), (8), (9)). We have focused our attention on the allowed region of the parameter space leading to a correct electroweak breaking, the fine-tuning problem, and the upper bounds on supersymmetric masses.

As a preliminary, we showed that some common approximations, such as considering only the top and stop contributions to ΔV_1 and/or disregarding the $\tilde{t}_L - \tilde{t}_R$ mixing, though acceptable for other purposes, lead to wrong results for $SU(2) \times U(1)$ breaking. In consequence, we have worked with the exact one-loop effective potential V_1.

Next, we have examined what region of the whole MSSM parameter space (i.e. the soft breaking terms $M_{1/2}, m_0, A, B$ plus μ and h_t) leads to a correct $SU(2) \times U(1)$ breaking, i.e. to the correct value of M_Z, a value of m_t consistent with the observations and no colour or electric charge breakdown. A comparison with the results of the "renormalization improved" tree-level potential V_0 [5] shows that the one-loop corrections enhance (and also displace) the allowed region of parameters. This, of course, is good news for the MSSM.

Our following step has been to analyse the top-fine-tuning problem. As was pointed out in ref. [5], h_t (the top Yukawa coupling) is the parameter to which M_Z is most sensitive. Using the ordinary criterion to avoid fine-tuning, i.e. $c \lesssim 10$ in the relation

$$\frac{\delta M_Z^2}{M_Z^2} = c\frac{\delta h_t^2}{h_t^2}\ , \tag{16}$$

strongly constrains the values of the MSSM parameters, leading to upper bounds on $M_{1/2}, m_0, \mu$, and thus on the masses of the new supersymmetric states (gluino, squarks, charginos, etc.). This analysis was performed in ref. [5] using the improved tree-level potential V_0. We find that the one-loop corrections substantially soften the degree of fine-tuning. This, again, is good news for the MSSM.

Finally, we have explored what the limitations of the ordinary criterion (16) are to parametrize the degree of fine-tuning. We comment on its ambiguities and show a type of (hypothetical) scenarios in which this criterion would be completely meaningless. Fortunately, this is not the case for the MSSM and, thus, the c parameter represents a sensible, but qualitative estimate of the degree of fine-tuning. A precise and non-ambiguous quantification of it can only be done once the supergravity breaking mechanism is known. In view of all this, we have conservatively relaxed the acceptable upper bound for c up to $c \lesssim 20$.

As a summary of the results the one-loop contributions (i) enhance (and displace) the allowed region of the MSSM parameters, (ii) soften the fine-tuning associated with the top quark (for large values of the MSSM parameters). These two facts, together with the fact that (iii) the upper bound on c should be conservatively relaxed, push up the upper bounds on the MSSM parameters obtained from the former V_0 analysis *and* the corresponding upper bounds on supersymmetric masses. This is reflected in table 1 for two "extreme" cases and in eq. (15). Our final conclusion is that the supersymmetric spectrum is not necessarily close to the present experimental limits, although the future accelerators (LHC, SSC) should bring it to light.

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Note added in proof

After sending this paper for publication we have received a paper by M. Olechowski and S. Pokorski [14], where the effect of the one-loop contributions to the electroweak breaking is also analyzed. Their results about the degree of fine-tuning are in agreement with ours.

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