

HU-SEFT R 1995-08
UCL-IPT-95-07
CPT-95/PE.3185

Top-Induced Electroweak Breaking in the Minimal Supersymmetric Standard Model

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Abstract

Severe constraints on parameters of the minimal supersymmetric standard model follow from a dynamical electroweak symmetry breaking mechanism dominated by top and stop loops. In particular, the lightest Higgs boson mass is expected to be smaller than 100 GeV.

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Introduction

The recent confirmation [1] of the experimental discovery of the top quark with a mass $m_t \approx 175\text{--}200$ GeV is yet another success of the standard model (SM) of electroweak interactions. Indeed, theoretical computations of radiative corrections to the SM confronted with the precise LEP data led to a range of values for m_t in remarkable agreement with the experimental results. To appreciate the significance of these theoretical and experimental achievements, it is worth recalling that calculations and measurements of the $g-2$ factor of the electron do not allow for a determination of the muon mass with anything even close to the precision of the top mass prediction in the SM !

Even though the SM is in remarkable agreement with all experimental facts, it is fair to say that the origin as well as the hierarchy of quark and lepton masses remain poorly understood. These issues are related to the spontaneous symmetry breaking (SSB) of the electroweak gauge symmetry. The usual tree-level Higgs mechanism for SSB is unsatisfactory in many ways and one of the most attractive alternatives suggested so far is due to Coleman and Weinberg [2] (see also [3]): in their approach, SSB is a dynamical effect induced by radiative corrections and not an ad-hoc input at the tree level of the theory.

Whatever the physics underlying SSB, the discovery of a Higgs scalar [4] is of course a tantalizing challenge to experimentalists. On the other hand it is clearly interesting and important to determine theoretically possible ranges of values i.e., bounds for the mass m_h of the Higgs in the SM and of the lightest scalar in as wide a class of models extending the SM as possible.

In the context of the SM viewed as an effective theory, i.e. valid up to some scale Λ , the usual Higgs mechanism leads to a lower bound on m_h determined from vacuum stability arguments [5], while an upper bound follows from the requirement that no Landau pole appears up to the scale Λ [6].

In a recent letter [7], very much inspired by the Coleman-Weinberg approach, it was suggested that the breaking of the electroweak gauge symmetry in the SM is a quantum dynamical effect driven by top quark loops. The physical motivation for this breaking mechanism comes from the fact that, with $m_t \simeq 200$ GeV, the Yukawa coupling of the top quark to the Higgs scalar is of order one and thus significantly larger than all other electroweak coupling constants. A straightforward consequence of this SSB mechanism is an upper bound on m_h of the order of 400 GeV. It is however tempting to speculate further that top loops are responsible for the physical effects of SSB, i.e. that the full “Higgs potential” is dynamically generated by quantum effects. The consistency of this point of view requires the effective scalar quartic interaction to be negligibly small if not zero, indicating a possible connection with the “triviality” of a $\lambda\phi^4$ theory [8]. Be that as it may, the phenomenological implementation of this additional assumption leads to a cut-off Λ of the order of 1 TeV and to a Higgs mass $m_h \approx 80$ GeV.

The purpose of the present paper is to explore a bit further the phenomenological implications of such a dynamical symmetry breaking mechanism in the context of the minimal supersymmetric extension of the standard model (MSSM) where the scalar quartic

interactions are functions of the electroweak couplings and, therefore, naturally small.

The attractive features of MSSM are well known [9]. Among others, it provides a plausible solution to some of the puzzles inherent to the SM like the “gauge hierarchy” problem. Furthermore, MSSM fits nicely with unification ideas since it is naturally embedded in many unified supergravity and superstring models. If there is new physics beyond the SM, MSSM is certainly one of the prime candidate theories to describe it.

In MSSM there are two $SU(2)_L \otimes U(1)$ Higgs doublets and, hence, in the end there will be three neutral Higgs scalars (two CP-even ones and one CP-odd one). For simplicity we will view MSSM as an effective theory with three characteristic scales : the cut-off Λ (which, as it will appear, plays almost no role in our arguments*), the global supersymmetry breaking scale M_S and, finally, the electroweak breaking scale given by a vacuum expectation value (VEV) v . Our line of argument starts as usual in the context of MSSM. First, supersymmetry is broken. As a result all SUSY partners of the SM particles become heavy with masses of the order of M_S . In the Higgs sector, the CP-odd scalar and one linear combination, characterized by a mixing angle β , of the two CP-even scalars become heavy with a mass again of the order of M_S , while the orthogonal CP-even combination remains light (with a positive mass). At this stage $SU(2)_L \times U(1)$ is still an exact gauge symmetry of the model.

We proceed by assuming that SSB of the gauge group is completely driven by top and stop loops since their couplings to the light scalar field are equal and large. Consistency of this picture of SSB requires the quartic scalar self coupling to be small. Putting it equal to zero leads to our final results

$$m_h \lesssim 100 \text{ GeV}, \quad M_S \lesssim 600 \text{ GeV} \quad \text{and} \quad 0.6 \lesssim \tan \beta \lesssim 1.6 .$$

The paper is organized as follows : in Section I we briefly review well known features of the Higgs sector in MSSM and in Section II we derive and present our results. We then conclude with some brief comments on these results as well as on earlier investigations concerning the Higgs masses in MSSM.

I The Higgs sector in MSSM

Let us recall [4] that in MSSM there are two complex scalar $SU(2)_L$ doublets of opposite hypercharge: $H_1 = \begin{pmatrix} H_1^0 \\ H_1^- \end{pmatrix}$ which couples to the charged leptons and to down-type quarks and $H_2 = \begin{pmatrix} H_2^+ \\ H_2^0 \end{pmatrix}$ which couples to up-type quarks only.

The tree level scalar potential for the real components of the CP-even neutral scalars

* Because of this weak dependence of the results on the cut off Λ , a renormalized version of the theory would lead to similar conclusions.

$h_1 \equiv \text{Re } H_1^0$ and $h_2 \equiv \text{Re } H_2^0$ reads

$$V_0 = \frac{1}{2}m_1^2 h_1^2 + \frac{1}{2}m_2^2 h_2^2 - m_{12}^2 h_1 h_2 + \frac{\tilde{g}^2}{32}(h_2^2 - h_1^2)^2, \quad (1)$$

where $\tilde{g}^2 = (g_1^2 + g_2^2)$ and $g_{1,2}$ are the usual $U(1)$ and $SU(2)$ gauge couplings.

Generically, the mass terms in Eq. (1) break supersymmetry softly. For our purposes we require the potential V_0 to be bounded from below, i.e.

$$m_1^2 + m_2^2 \geq 2m_{12}^2$$

and to be such that the gauge symmetry $SU(2)_L \otimes U(1)$ remains unbroken at the tree level:

$$m_1^2 m_2^2 > m_{12}^4.$$

As usual in this context, we assume that all SUSY partners of SM particles have become heavy with masses of the order of the global SUSY breaking scale M_S and that one linear combination of the neutral scalars

$$h = h_1 \cos \beta + h_2 \sin \beta \quad (2)$$

remains light, while the orthogonal combination

$$H = -h_1 \sin \beta + h_2 \cos \beta \quad (3)$$

has a mass of the order of M_S .

The relevant part of the tree level potential for the field h then reads

$$V_0 = \frac{1}{2}\mu^2 h^2 + \frac{1}{32}\tilde{g}^2 \cos^2 2\beta h^4 \quad (4)$$

with $\mu^2 > 0$ and the mixing angle β , defined by Eqs. (2)-(3) is determined by the mass parameters in Eq. (1).

We insist on the fact that at this stage only supersymmetry has been broken: h and H are normal scalar fields with positive masses and no vacuum expectation values. We notice that many supergravity inspired models imply the boundary condition $m_1^2 = m_2^2$ with $m_{12}^2 \neq 0$ around the GUT scale. In this case $m_H^2 \approx m_1^2 + m_{12}^2$, $\mu^2 \approx m_2^2 - m_{12}^2$, and $\beta = \pm\pi/4$. It is interesting to note that the latter relation $\beta = \pm\pi/4$ would also emerge from the “triviality” of the scalar potential V_0 , i.e. from the requirement $\partial^4 V_0 / \partial h^4 = 0$.

In Eq. (4) we have omitted terms of the form $h^2 H^2$, $h H^3$ and $h^3 H$. It turns out that the electroweak gauge boson contributions to the one loop effective potential are negligible because of the relative smallness of the gauge couplings compared to the top Yukawa coupling. We expect the same suppression to occur for H virtual contributions since its couplings to h are also given by the gauge couplings. In addition m_H could also be assumed to be much larger than M_S .

II The effective potential at one loop

We now proceed in a completely standard fashion with the computation of the dominant one-loop radiative corrections V_1 to the scalar potential V_0 . In other words we compute the infinite series of diagrams given in Fig. 1a with massless top quarks in the loop, and in Fig. 1b with left and right stops in the loop. We assume that the stops have the same soft mass m_s (of the order of M_S) and neglect possible left-right mixing effects**.

The resulting effective potential is well known and reads

$$V = V_0 + V_1$$

with V_0 given by Eq. (4) and

$$V_1 = \frac{N_c}{8\pi^2} \int_0^{\Lambda^2} dk^2 k^2 \ln \left(\frac{k^2 + m_s^2 + g_t^2 h^2/2}{k^2 + g_t^2 h^2/2} \frac{k^2}{k^2 + m_s^2} \right), \quad (5)$$

N_c is the number of colours, Λ is the cut-off scale and $g_t = h_t \sin \beta$ is the top coupling to the scalar h ; h_t is the supersymmetric Yukawa coupling of the top to H_2 and the factor $\sin \beta$ follows from Eq. (2).

Note that when $m_s = 0$, $V_1 = 0$ as it should, while for $m_s \rightarrow \infty$ we recover the top loop radiative corrections to the scalar potential in an effective SM.

Minimization of the potential V now yields a non-trivial minimum $\langle h \rangle \equiv v \neq 0$: top and stop loops effects have thus induced the SSB of the electroweak gauge group. At $\langle h \rangle = v$, one obtains the equation

$$\mu^2 + \frac{1}{2} m_Z^2 \cos^2 2\beta + \frac{N_c}{4\pi^2} \frac{m_t^2}{v^2} \left\{ m_t^2 \ln \left(1 + \frac{\Lambda^2}{m_t^2} \right) - (m_t^2 + m_s^2) \ln \left(1 + \frac{\Lambda^2}{m_t^2 + m_s^2} \right) \right\} = 0, \quad (6)$$

where

$$m_Z^2 = \frac{1}{4} \tilde{g}^2 v^2, \quad (v = 246 \text{ GeV}) \quad (7a)$$

and

$$m_t^2 = \frac{1}{2} g_t^2 v^2. \quad (7b)$$

Since μ^2 is assumed positive, the expression in brackets in Eq. (6) must be negative. This is always the case for any $m_s^2 > 0$. We postpone further comments on Eq.(6) to the next section.

For now, let us compute the lightest Higgs boson mass at the scale v i.e. $m_h^2 = (\partial^2 V / \partial h^2)|_{h=v}$. Using Eq. (6), this yields the result

$$m_h^2 = m_Z^2 \cos^2 2\beta + \frac{N_c}{2\pi^2} \frac{m_t^4}{v^2} \left\{ \ln \left[\left(1 + \frac{m_s^2}{m_t^2} \right) \left(\frac{\Lambda^2 + m_t^2}{\Lambda^2 + m_t^2 + m_s^2} \right) \right] \right.$$

**Notice that after the breaking of the gauge symmetry when the top quark has acquired a mass m_t , the mass of both left and right stops will be given by $m_t^2 \simeq m_s^2 + m_t^2$, when their mixing is neglected.

$$\left. - \frac{\Lambda^2 m_s^2}{(\Lambda^2 + m_t^2)(\Lambda^2 + m_t^2 + m_s^2)} \right\} . \quad (8)$$

Notice that m_h^2 is weakly cut-off dependent for $\Lambda \gtrsim 1$ TeV and remains finite in the limit $\Lambda \rightarrow \infty$.

We thus obtain the upper bound

$$m_h^2 \leq m_Z^2 \cos^2 2\beta + \frac{N_c}{2\pi^2} \frac{m_t^4}{v^2} \ln \left(1 + \frac{m_s^2}{m_t^2} \right) . \quad (9)$$

Eqs. (8) and (9) are our first results. They follow directly from the assumption that top and stops loops dominate the radiative corrections to the scalar potential and drive the SSB of the gauge group. In Fig. 2 we plot the upper limit on the light Higgs mass as given by Eq. (9) for two extreme values of the mixing angle β .

For our physical assumptions to be coherent one should argue further that the h scalar self interactions remain negligible, at least at the scale v . This is easily implemented phenomenologically by the condition

$$\frac{\partial^4 V}{\partial h^4} \Big|_{h=v} \simeq 0 . \quad (10)$$

Taken as an exact equation, Eq. (10) looks like a “triviality condition” on a scalar field theory. Again we postpone comments on this point to the next section but for the moment we consider Eq. (10) as a phenomenological constraint on the effective potential which is important to guarantee, a posteriori, that h scalar loops will not affect the SSB mechanism induced by top and stop loops. Putting it differently, if the physics of SSB is indeed dominated by radiative corrections due to top and stop loops then scalar self interactions must be small, at least at the scale where SSB occurs.

After some simple algebra, Eq. (10) takes the form

$$\frac{1}{4} \tilde{g}^2 \cos^2 2\beta + \frac{N_c}{2\pi^2} \frac{m_t^4}{v^4} \{f(x, y) - f(0, y)\} \cong 0 , \quad (11)$$

where

$$x = \frac{m_s^2}{m_t^2} , \quad y = \frac{\Lambda^2}{m_t^2} \quad (12)$$

and

$$\begin{aligned} f(x, y) = & -\ln\left(1 + \frac{y}{1+x}\right) - \frac{9+x}{1+x+y} + \frac{4(2+x)}{(1+x+y)^2} \\ & - \frac{8}{3} \frac{1+x}{(1+x+y)^3} - \frac{4x(1+2x)}{3(1+x)^2} . \end{aligned} \quad (13)$$

We now proceed to discuss the phenomenological content of Eqs. (8) and (11).

For simplicity let us start with the limiting case (of large cut off) $y \rightarrow \infty$. For Eq.(11) to admit a solution we must then demand that

$$\lim_{y \rightarrow \infty} [f(x, y) - f(0, y)] \lesssim 0$$

or, explicitly, that

$$F(x) \equiv \ln(1+x) - \frac{4x(1+2x)}{3(1+x)^2} \lesssim 0. \quad (14)$$

The function $F(x)$ is plotted in Fig. 3. The equation $F(x) = 0$ has a unique solution, namely $x_0 = 8.65$, which is the maximum value of x allowed by Eq. (11). This requires

$$m_s \lesssim 2.9 m_t. \quad (15)$$

On the other hand, the function $F(x)$ reaches its minimum value for $x \simeq 2.16$ and $F(2.16) \simeq -0.38$. This immediately leads to an upper bound for $\cos^2 2\beta$ or to a range of values for $\tan \beta$ given by

$$\left(\frac{1 - 0.32g_t^2}{1 + 0.32g_t^2} \right)^{1/2} \lesssim \tan \beta \lesssim \left(\frac{1 + 0.32g_t^2}{1 - 0.32g_t^2} \right)^{1/2}. \quad (16)$$

The solution of Eq. (11) in the limit $\Lambda \rightarrow \infty$ is expressed in Fig. 4 as a plot of $\tan \beta$ in terms of m_s for the central values of m_t in the two Fermilab experiments, namely $m_t = 176$ GeV and $m_t = 199$ GeV.

Eq. (11) thus provides two important constraints on the parameters of MSSM : an upper limit on m_s (Eq. (15)) and a range of values for $\tan \beta$ (Eq. (16)).

Provided $\tan \beta \neq 1$, we can solve Eq. (11) for $\cos^2 2\beta$ and substitute the result into Eq. (8) to obtain

$$m_h = \frac{2}{\pi} \frac{m_t^2}{v} \left\{ \frac{y^3}{(1+y)^3} - \frac{1}{2} \frac{(3x+2y+3x^2+3xy)y^2}{(1+x)^2(1+x+y)^3} \right\}^{1/2}. \quad (17)$$

Notice that in the limit $x \rightarrow \infty$, Eq. (17) reproduces the result found in Ref. [7], namely

$$m_h = \frac{2}{\pi} \frac{m_t^2}{v} \left\{ 1 + O\left(\frac{1}{y}\right) \right\}. \quad (18)$$

In the limit $y \rightarrow \infty$ Eq. (17) becomes

$$m_h = \frac{2}{\pi} \frac{m_t^2}{v} \left\{ 1 - \frac{1}{2} \frac{(2+3x)}{(1+x)^2} \right\}^{1/2}, \quad (19)$$

meaning, in particular, that m_s is a natural cut off for the effective standard model considered in Ref. [7] if $x \gg 1$.

Eq. (19) thus gives an upper bound on m_h in the limit $\Lambda \rightarrow \infty$. To illustrate the weak sensitivity of our results on the cut-off scale Λ , we have plotted in Fig. 5 m_h (as given by Eq. (17)) as a function of m_s for $\Lambda = 1$ TeV and $\Lambda = \infty$.

Finally, let us briefly comment on our results in the particular case $\beta = \pm\pi/4$. The Higgs mass is then given by the second term in Eq. (8), namely

$$m_h^2 = \frac{N_c}{2\pi^2} \frac{m_t^4}{v^2} \left\{ \ln \frac{(1+x)(1+y)}{1+x+y} - \frac{xy}{(1+y)(1+x+y)} \right\}, \quad (20)$$

where x and y now satisfy the equation $f(x, y) = f(0, y)$ with $f(x, y)$ as in Eq. (13). For $\Lambda \rightarrow \infty$ and $m_s \simeq 2.9 m_t$, one finds $m_h \simeq 74$ GeV (resp. 94 GeV) for $m_t = 176$ GeV (resp. 199 GeV), as illustrated in Fig. 5. This concludes the presentation of our results. Clearly the most striking phenomenological implications of our model for SSB of the electroweak gauge group in the context of MSSM are the ones already mentioned in the introduction, namely

$$m_h \lesssim 100 \text{ GeV} , \quad (21)$$

$$m_s \lesssim 600 \text{ GeV} \quad (22)$$

and

$$0.6 \lesssim \tan \beta \lesssim 1.6 . \quad (23)$$

III Comments

The basic physical idea of the present paper is that in MSSM, spontaneous symmetry breaking of the electroweak gauge group is a quantum effect predominantly driven by top and stop loops. Implementation of this SSB scenario puts very severe constraints on the allowed range of values for some of the most important phenomenological parameters in MSSM. These constraints are summarized in Eqs. (21)-(23) and imply in particular, a very light Higgs scalar which should be observable [10] at LEP 200 if not before and a rather low supersymmetry breaking scale.

To conclude, we present a series of miscellaneous comments which among others express our views on the reliability of the calculations presented here as well as on the relation of our work to other approaches to SSB in MSSM.

1. In the effective potential, we have neglected W and Z boson loops as well as the corresponding gauginos loops. An explicit calculation of their effect, say on the Higgs boson mass, shows indeed that it is at most of the order of one percent. The main reason for this is of course that the gauge couplings g_1 and g_2 are significantly smaller than g_t - mutatis mutandis we expect the same argument to hold for H and the corresponding higgsino loops. Alternatively, one can always push the mass of H and its SUSY partner $\tilde{H}$ high enough so as to make their effect on V_1 completely negligible.

2. It should be noted that β is originally a mixing angle defined by the diagonalization of a mass matrix (Eq. (1)). At the one loop level, the same β (with small corrections) is related to the vacuum expectation values $v_{1,2}$ of the fields $H_{1,2}^0$. Indeed, since $\langle H \rangle = 0$ we obtain $\tan \beta = v_2/v_1$, as usual. The VEV of h is of course given by $v = (v_1^2 + v_2^2)^{1/2}$.

3. On several occasions we have referred to the “triviality issue” in a $\lambda\phi^4$ theory. Let us recall that in a pure $\lambda\phi^4$ theory the renormalized coupling constant λ_R must be zero or else the theory blows up at some scale (related to the value of λ). The argument is essentially non perturbative [8]. We are not suggesting that triviality should be imposed order by order in the loop expansion of the effective potential although such a procedure probably deserves further study. For the purposes of this paper, it is quite sufficient to

take the more phenomenological point of view emphasized in section II : to insure that top and stop loops do indeed dominate the effective potential the quartic scalar self coupling must be small, at least at the scale v .

4. The results presented in section II are weakly cut-off dependent. For a very large cut off we expect QCD corrections to our results to become huge. For Λ around 1 TeV on the other hand, QCD corrections remain small.

5. Two sine qua non conditions for our approach to SSB to make sense is that higher loop corrections to the one loop effective potential remain small and that the “one loop vacuum” is stable. As usual it is easy to argue that higher loop effects will indeed be negligible [2] but vacuum stability is a much more difficult condition to analyze. We have no insight on this problem.

6. Of course many interesting papers have already been published on Higgs mass predictions or bounds in MSSM. They can be roughly classified according to the value of μ^2 : in the usual scenario μ^2 is supposed to be negative at the weak scale [11] while the critical case $\mu^2 = 0$ has also been considered [12]. In both approaches, it has been emphasized that a heavy top leads to large corrections to the tree level bound or prediction for m_h .

In the case $\mu^2 < 0$, m_h is expected to be larger than 120 GeV while for $\mu^2 = 0$, m_h is around 100 GeV. It is thus not surprising that in a scenario with $\mu^2 > 0$, m_h is predicted to be even lighter.

Acknowledgements

We are grateful to K. Huitu and A. Kupiainen for useful discussions. One of us (P.C.) would like to thank the Research Institute for High Energy Physics (SEFT), University of Helsinki for the hospitality. This work was partially supported by the EEC Science Project SC1-CT91-0729.

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Figure Captions

Fig. 1 One-loop radiative corrections to the scalar potential V_0 due to top quark and stop contributions. The infinite series of diagrams in Fig. 1a represents the massless top quark contribution to the one-loop effective potential, while the one in Fig. 1b corresponds to left and right stops with soft mass m_s .

Fig. 2 Upper limit on the Higgs mass (Eq. (9)) under the assumption that top and stops loops dominate the radiative corrections to the scalar potential and drive the SSB of the gauge group. The plot is given for the extreme values of the mixing angle $\beta = \pi/2$ and $\beta = \pi/4$ and for $m_t < 200$ GeV.

Fig. 3 Function $F(x)$ as given by Eq. (14) ($x = m_s^2/m_t^2$). The solution of the equation $F(x) = 0$ gives an upper bound on the soft mass parameter m_s , $m_s \lesssim 2.9 m_t$.

Fig. 4 Allowed range of values of $\tan \beta$ as a function of m_s in the limit $\Lambda \rightarrow \infty$ for the central values of m_t reported by the two Fermilab experiments, namely $m_t = 176$ GeV and $m_t = 199$ GeV.

Fig. 5 Higgs boson mass m_h as a function of m_s for $m_t = 176$ GeV and $m_t = 199$ GeV. The curves are plotted for $\Lambda = 1$ TeV (dashed lines) and $\Lambda \rightarrow \infty$ (solid lines), which indicate the weak dependence of the Higgs mass on the cut off. The maximal values for m_h correspond to $\tan \beta = 1$.

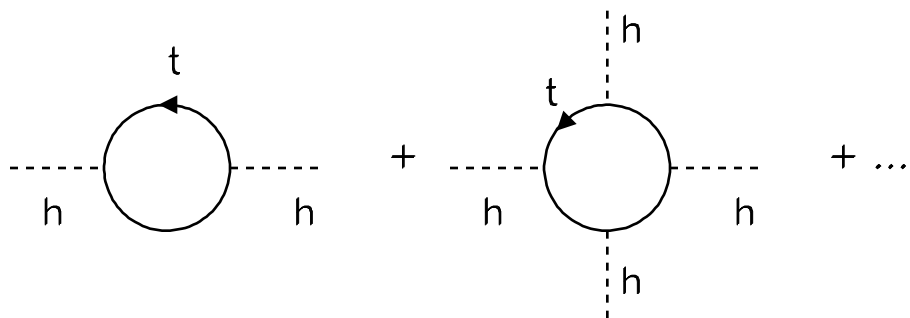


Fig. 1a

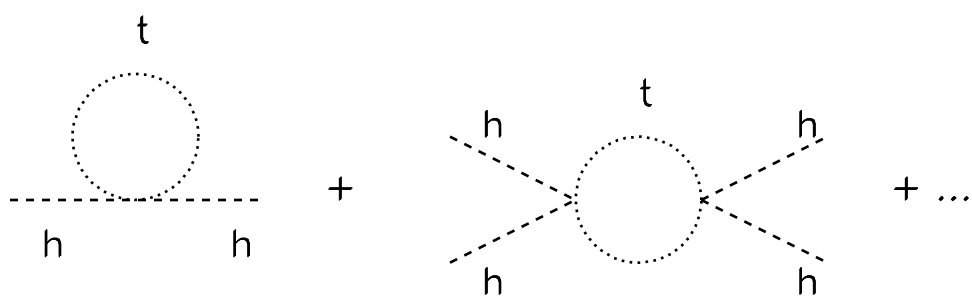


Fig. 1b

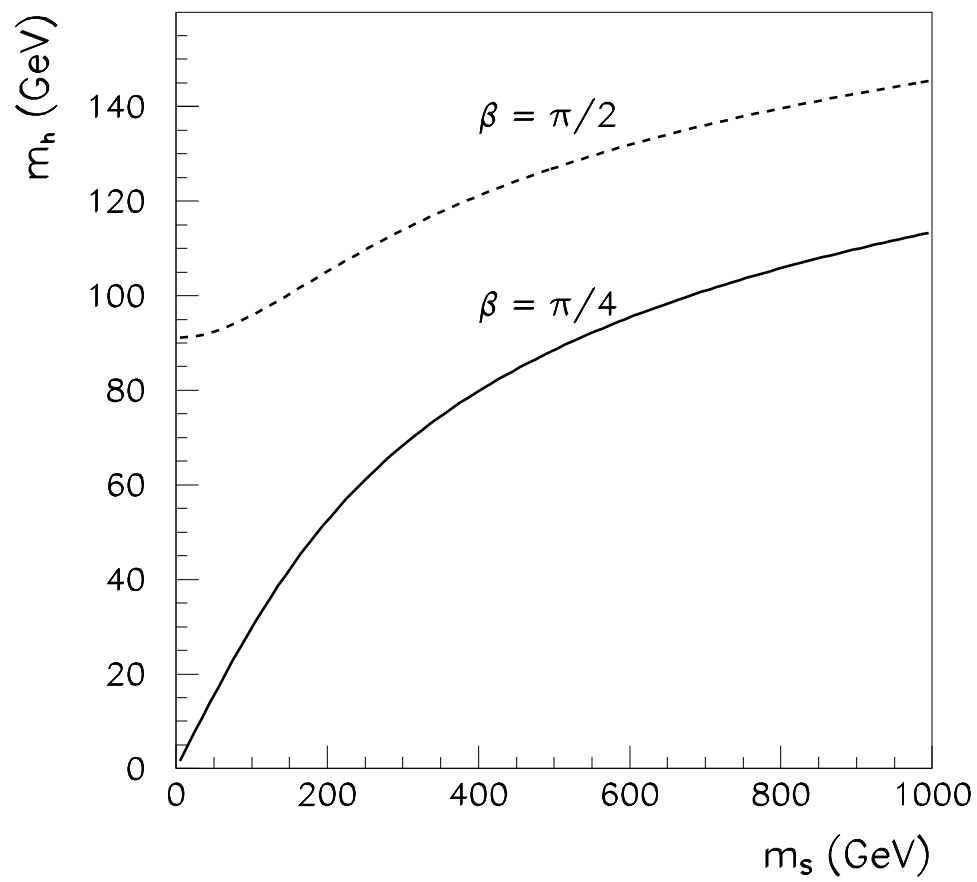


Fig. 2

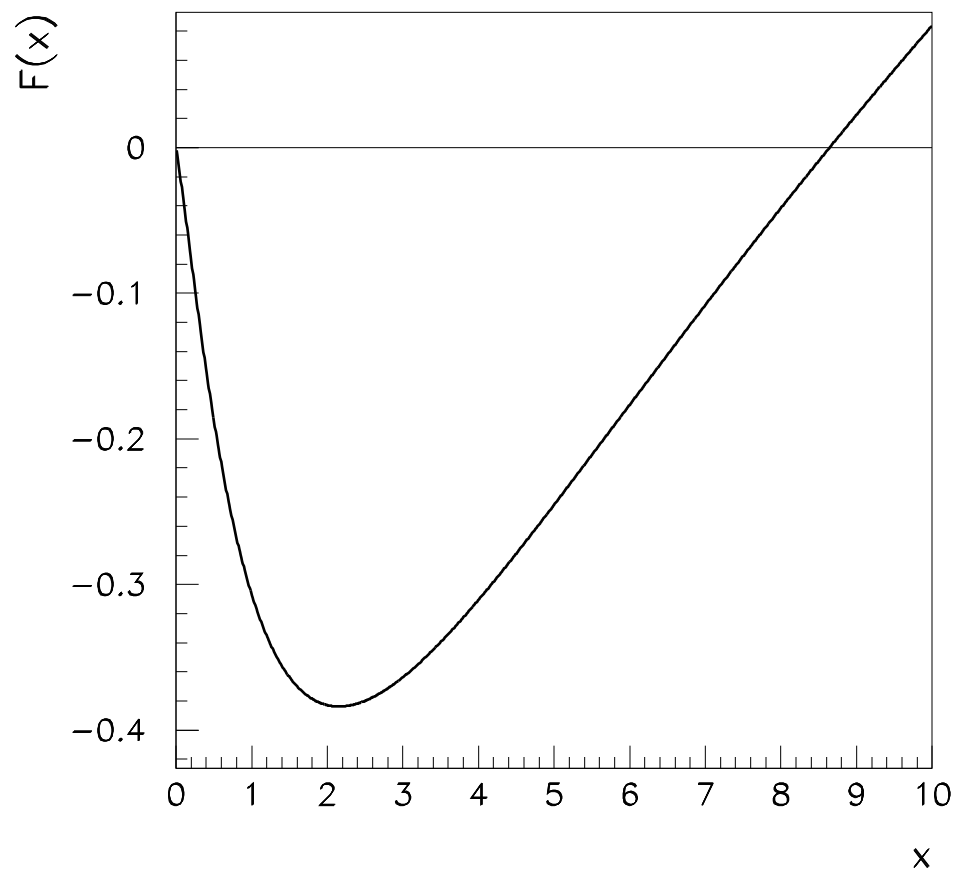


Fig. 3

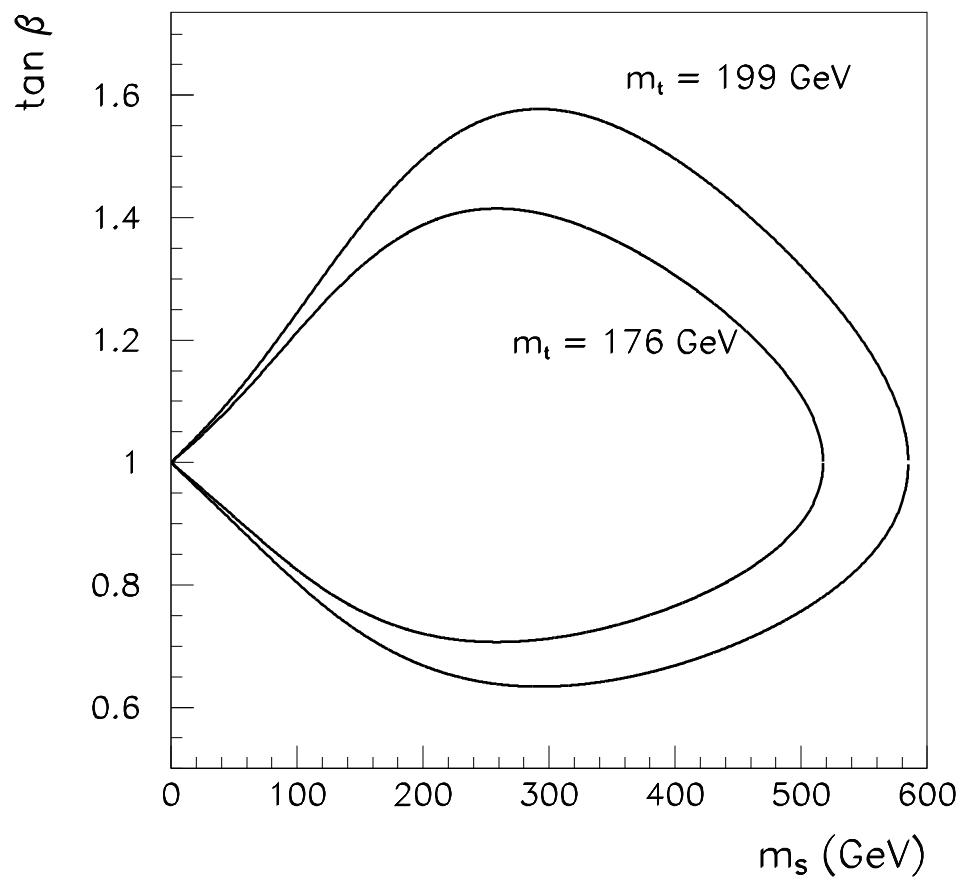


Fig. 4

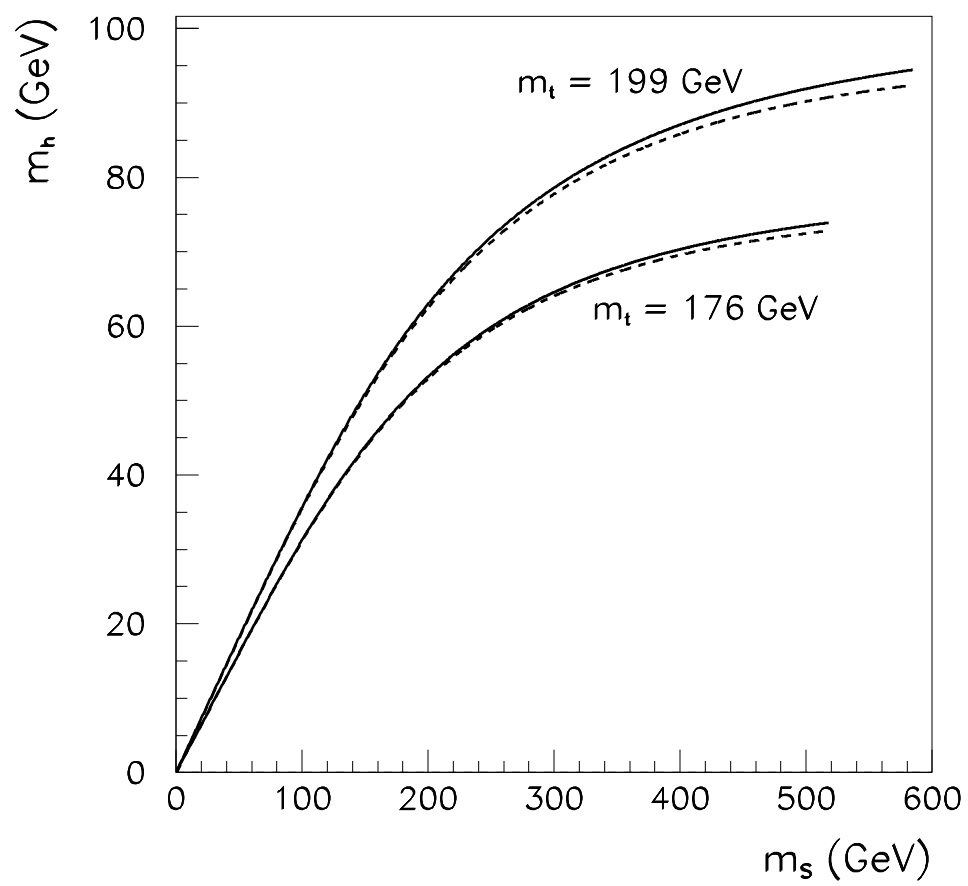


Fig. 5