Constraints on New Physics from the 

Higgs and Top Masses

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Abstract

Triviality and vacuum stability bounds on the Higgs and top quark masses in a rather general class of supersymmetric extensions of the Standard Model are compared with the corresponding bounds without supersymmetry. Due to generic differences of those bounds we find that experimental knowledge of the Higgs and top masses may provide a “pointer” into one of these directions. Depending on the values of the masses, however, both scenarios or none could also be allowed.

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The Standard Model of electro–weak interactions is a quantum field theory which is in agreement with all existing experimental data. This includes also some evidence for radiative corrections as required by the theory. Nevertheless it is for different reasons very likely that the Standard Model is embedded into a larger framework. One of the most important reasons is the so called hierarchy problem which is based on the observation that the quadratic divergences of the Higgs sector make it hard to explain a big hierarchy between $v \approx 175 \text{ GeV}$ and a very high scale of new physics $\Lambda$. The hierarchy problem is, however, only a strong argument for new physics beyond the Standard Model if the cutoff has a physical meaning. In the renormalizable Standard Model itself the problem does not exist since it is absorbed by renormalization. One might therefore take an extreme attitude and dismiss all those arguments for new physics.

But even then the ad hoc invention of the Higgs sector in order to break the electro–weak symmetry does not necessarily imply that fundamental scalar fields must exist. Like in the case of the Ginzburg–Landau description of superconductivity these scalars might turn out to be just an effective parametrization of some more complex dynamical scenario. However, independently of the question whether the Standard Model is just an effective field theory up to some scale $\Lambda$ the allowed range of parameters is restricted. These restrictions stem from the possibility that the vacuum of the theory can be unstable \cite{1} or that the model is “trivial”, which means that the only consistent version of the theory is the free, non–interacting case \cite{2}. In the language of running coupling constants these two problems can be phrased as the possibility that the Higgs self coupling $\lambda(\mu)$ becomes negative such that the Higgs potential is unbounded from below, or the possibility that one of the running couplings develops a Landau singularity \cite{3}. Both type of problems can in principle occur at an arbitrarily high scale $\mu$, but in order to be physically relevant one has to require that $\mu < \Lambda$, where $\Lambda$ is the range of validity of the Standard Model\cite{1}. If the hierarchy problem is not solved in some unexpected way $\Lambda$ should probably not exceed a few $\text{TeV}$. Apart from the Higgs self coupling $\lambda$ the other possibly large coupling in the Standard Model is the top quark Yukawa coupling. Accordingly these restrictions lead to constraints on the physical Higgs and top quark masses \cite{4–7}.

Alternatively, if fundamental scalars really exist, a natural solution to the hierarchy problem is given by supersymmetry. This is because scalars emerge naturally and quadratic divergences are canceled beyond the supersymmetry breaking scale $\Delta$, thus the hierarchy problem is solved if $\Delta \approx 1 \text{ TeV}$. The supersymmetric extension of the Standard Model is, however, by no means unique. But it is very natural to assume, that any supersymmetric

\footnote{Note, however, that a Landau singularity at the embedding scale can be considered as an indication of compositness at this scale \cite{8, 9}.}
extension of the Standard Model is a consistent field theory up to a GUT or even the Planck scale; after all this possibility is the main motivation for the introduction of supersymmetry. This implies again the absence of Landau singularities for the running couplings, now up to these very large scales. The couplings under consideration are Higgs self couplings and the top quark Yukawa coupling as before; hence one obtains again constraints on the physical Higgs and top quark masses. Within a general supersymmetric extension, however, lower bounds on the lightest Higgs mass from the condition of vacuum stability cannot be obtained due to the different form of the scalar potential and the radiative corrections.

Bounds on the mass of the lightest Higgs scalar in the framework of the so-called minimal extension have been discussed in much detail recently [10–18]. Apart from refs. [11, 14], however, constraints from a consistent “high energy input” have not been taken into consideration in these investigations. Within non–minimal extensions as, e.g., the addition of a gauge singlet to the Higgs sector [19–28] these bounds become weaker. Including the leading log radiative corrections the corresponding upper bounds have recently been computed in [26–28]. This latter model can actually be viewed as the appropriate testing ground for the general assumption of supersymmetry. It is sufficiently general and contains the minimal extension for special choices of its parameters. The addition of further doublets to the Higgs sector would not change the upper bound on the mass of lightest Higgs field [22, 28, 25].

A comparison of the constraints on the Higgs and top quark masses within the Standard Model and its minimal supersymmetric extension has recently been performed in [30]. There, however, the Standard Model was assumed to remain valid up to scales beyond 10^{10} GeV, and just the minimal supersymmetric extension was considered. Also the triviality constraint on the top quark Yukawa coupling was not implemented. In contrast we will use the non–minimal extension described above, which allows a more general supersymmetric scenario. Furthermore we believe that in the absence of supersymmetry it is sensible to require the absence of Landau singularities only up to a few TeV, since the unsolved hierarchy problem will very likely require such a low embedding scale.

Below we will sketch the derivation of the constraints on the Higgs and top quark masses for the two cases beyond the leading log approximation, where we make use of results obtained already elsewhere. From a comparison of these constraints we can learn, once the Higgs and top masses are experimentally known (or better constraint), whether perturbative supersymmetry up to 10^{16} GeV is allowed or excluded, or whether an unspecified embedding of the Standard Model at a few TeV (or higher) is allowed or excluded. We

\footnote{There might, however, exist solutions of the hierarchy problem within the Standard Model which would then require to take the model serious up to the GUT– or even the Planck scale [31].}
will discuss whether some experimental regions can be understood as pointers into one of those directions.

Within the Standard Model, the two undetermined couplings $g_t$ and $\lambda$, which are related to the unknown top and Higgs masses via $g_t = m_t/v$ and $\lambda = m_H^2/2v^2$, can develop Landau singularities or an unstable potential even at rather low scales. The renormalization group flow is given by $dg_t/dt = \beta_t$ and $d\lambda/dt = \beta_\lambda$ where

$$16\pi^2 \beta_t = \left( \frac{9}{2} g_t^2 - \frac{17}{12} g_1^2 - \frac{9}{4} g_2^2 - 8 g_3^2 \right) g_t ,$$

and

$$16\pi^2 \beta_\lambda = \left( 12 \lambda_2 - (A - 12 g_t^2) \lambda + B - 12 g_t^4 \right).$$

Here $t = \ln (\mu/\mu_0)$ and

$$A = 3 g_1^2 + 3 g_2^2 ; \quad B = \frac{3}{4} g_1^4 + \frac{3}{2} g_1^2 g_2^2 + \frac{9}{4} g_2^4 .$$

From the above beta functions we can immediately read off three possible problems:

- If $g_t$ is large the running coupling $g_t(\mu)$ can develop a Landau pole. For large $g_t$ eq. (4) can be approximated by $16\pi^2 \beta_t = 9/2 g_t^2$, which leads after integration to the approximate solution

$$\frac{1}{g_t^2(\mu)} = \frac{1}{g_t^2(\mu_0)} - \frac{9}{16\pi^2} \ln \left( \frac{\mu}{\mu_0} \right) .$$

The appearance of a Landau pole in $g_t(\mu)$ (i.e. a zero in $1/g_t^2(\mu)$) in the physical region below the embedding scale $\Lambda$ is avoided if the top mass is limited by

$$\frac{m_t^2}{v^2} = g_t^2(m_t) < \frac{16\pi^2}{9 \ln (\Lambda/m_t)} \simeq \frac{16\pi^2}{9 \ln (\Lambda/v)} ,$$

which leads to a bound which is in the Standard Model weaker than the other two bounds below. The approximation above describes the true result for small $\Lambda$ actually quite well. For large $\Lambda$ the bound (5) is too stringent which is immediately clear from the omission of the gauge couplings in the $\beta$–function. The full $\Lambda$–dependence of this bound with all running gauge couplings taken into account was discussed in [5].

- For large $m_H$, i.e. large $\lambda$ and small $g_t$ the $\beta$–function eq. (2) simplifies and becomes $16\pi^2 \beta_\lambda \simeq 12\lambda^2$. Integration leads then to

$$\frac{1}{\lambda(\mu)} = \frac{1}{\lambda(\mu_0)} - \frac{3}{4\pi^2} \ln \left( \frac{\mu}{\mu_0} \right) .$$
To avoid a Landau pole of $\lambda(\mu)$ in the physical region one must require
\[
\frac{m_H^2}{2v^2} = \lambda(m_H) < \frac{4\pi^2}{3\ln(\Lambda/m_H)} \simeq \frac{4\pi^2}{3\ln(\Lambda/v)}.
\] (7)

This approximate “triviality” bound for $m_H$ is again quite accurate for small $\Lambda$ while it is somewhat too stringent for large $\Lambda$. The full problem has been studied in detail with all effects included in ref. [5]. Note that this full result has also a weak top mass dependence.

- Finally for small $\lambda$ (and moderate $g_t$) the $\beta$–function eq. (3) can be simplified to become $16\pi^2\beta_\lambda \simeq B - 12g_t^4$. This leads to the approximate solution
\[
\lambda(\mu) = \lambda(\mu_0) + \frac{B - 12g_t^4}{16\pi^2} \ln \left( \frac{\mu}{\mu_0} \right).
\] (8)

From eq. (8) one can infer immediately that the solution can turn negative for $12g_t^4 > B$ which would change the sign of the quartic coupling leading to an unbounded potential. This must be avoided in the physical region below $\Lambda$.

Eq. (8) together with $\lambda(\Lambda) > 0$ and $\lambda(\mu_0) = m_H^2/2v^2$ translates into a lower bound on $m_H$ for large $m_t$. The approximation leads to
\[
\frac{m_H^2}{2v^2} > \frac{12m_t^4 - Bv^4}{16\pi^2v^4} \ln \left( \frac{\Lambda}{m_H} \right),
\] (9)

which shows how the bound starts at a certain value of $m_t$, and how it grows with $\Lambda$. The bound (9) is, however, typically somewhat above the full numerical result [5]. A detailed numerical study of eq. (2) with a number of other effects included (such as newer data, the most important two loop contributions to the $\beta$–functions, thresholds etc.) was performed in ref. [6]. Note that $\lambda(\mu)$ in eq. (8) can become negative immediately for $\mu$ above $m_H$ if the initial value of $\lambda(m_H)$ goes to zero and if $m_t$ is big enough to change the sign of the $\beta$–function. This explains the $\Lambda$ independence of this bound for very small Higgs masses.

When the three bounds discussed above are combined, we see that the allowed region in the Higgs–top mass plane is bounded to a $\Lambda$ dependent range around the origin (see Fig. 2 in [5]). Since the development of Landau pole(s) and of an unstable vacuum can be understood as “accidents” of the renormalization group flow it is also intuitively clear why the bounds are most restrictive for highest $\Lambda$, i.e. the largest running distance. Even though $\Lambda$ is in principle a free parameter we know that, for large $\Lambda$, the hierarchy problem will reappear as soon as we actually specify an embedding of the Standard Model. Unless some unusual
mechanism solves the hierarchy problem within the Standard Model this implies probably that \( \Lambda \) should not be very large, most likely only a few \( TeV \). In that case the bounds become weak, but they are still very interesting\(^3\). We will include the precise numerical results for the bounds just discussed in Fig. \[\] in the comparison at the end.

As outlined in the introduction, we will also discuss bounds on the Higgs and top masses within a non–minimal supersymmetric extension of the Standard Model. The Higgs sector of this non–minimal extension consists of two Higgs doublets, \( H_1, H_2 \) and a singlet \( S \). It is also motivated by the fact that it can get along with dimensionless supersymmetric couplings (no \( \mu H_1H_2 \) term in the superpotential), so that the electro–weak scale is introduced through the soft breaking terms only. (Possible additional dimensionful couplings will not modify the considerations below.) Since it is more general than the minimal supersymmetric extension, it is less restrictive; in particular already the tree level upper bound on the mass of the lightest Higgs scalar is not given by \( M_Z \), but depends – in some analogy to the non–supersymmetric model – on a dimensionless coupling \( \lambda \) \[^\[12\]\].

All relevant dimensionless couplings appear in the superpotential in the form

\[
W = g_t Q_L H_2 T_R + \lambda H_1 H_2 S + \frac{\kappa}{3} S^3. \tag{10}
\]

Here \( Q_L \) denotes the doublet containing the left–handed top and bottom quarks, \( T_R \) the right–handed top quark, and the vacuum expectation value \( v_2 \) of the Higgs doublet \( H_2 \) generates a top quark mass

\[
m_t = g_t v_2. \tag{11}
\]

In addition we take the following soft supersymmetry breaking trilinear couplings and masses into account:

\[
V_{soft} = (A_t g_t Q_{t,L} H_2 T_R + A \lambda H_1 H_2 S) + h.c. + m_1^2 |H_1|^2 + m_2^2 |H_2|^2 + m_S^2 |S|^2. \tag{12}
\]

Additional terms play no role subsequently. For the derivation of an upper bound of the lightest Higgs scalar of the model we adopt the following strategy: We consider the 2 by 2 mass matrix of the scalar neutral \( H_1 - H_2 \) sector and study its lightest eigenvalue, which constitutes such an upper bound. In this mass matrix we include the leading radiative corrections induced by top-quark and top-squark loops. Here we neglect the bottom quark mass and a possible splitting between the top squarks. The contributions of the gauge and the Higgs sector have been found to affect the final result only by \( \sim 5 \, GeV \) \[^\[12, 17\]\] in the direction of decreasing the upper bound on \( m_H \). (Also in the case of the extended Higgs

\[^3\]Note, however, that it would still be interesting if these bounds were violated experimentally for some larger \( \Lambda \) since this would establish an experimental upper limit on the range of validity of the Standard Model.
sector by the singlet these contributions can be estimated to be numerically unimportant.)

Two loop effects have been found to be of the order of \( \sim 5 \text{ GeV} \) \(^{16}\) and the difference between the pole mass and the second derivative of the effective potential \( \sim 3 \text{ GeV} \) \(^{17, 18}\). Hence we are on the safe side if we add 10 GeV to our upper bound on \( m_H \) obtained below in order to take these contributions into account. Furthermore we neglect terms of \( O(M_Z^2/m_t^2) \), which point into the lower direction anyway \(^{12}\). Now, with the help of the results of \(^{11}\), one finds the following elements of the mass matrix \( M_{ij} \) after elimination of \( m_2^2 \) and \( m_1^2 \) by means of the extremal equations:

\[
M_{11} = M_Z^2 \cos^2(\beta) + \Delta \tan(\beta) - \frac{3m_t^4 \lambda^2}{8\pi^2 v_2^2 m_{sq}^4} [A_t + \lambda \cot(\beta)]^2 ,
\]

\[
M_{22} = M_Z^2 \sin^2(\beta) + \Delta \cot(\beta) + \frac{3m_t^4}{8\pi^2 v_2^2 m_{sq}^2} \left[ 2 \ln \left( \frac{m_{sq}^2}{m_t^2} \right) + \frac{2A_t(A_t + \lambda \cot(\beta))}{m_{sq}^2} - \frac{A_t^2(A_t + \lambda \cot(\beta))^2}{6m_{sq}^4} \right] ,
\]

\[
M_{12} = M_{21} = -M_Z^2 \sin(\beta) \cos(\beta) - \Delta + \lambda^2 w \sin(2\beta) + \frac{3m_t^2 \lambda}{8\pi^2 v_2^2 m_{sq}^2} \left[ 1 - \frac{A_t(A_t + \lambda \cot(\beta))}{6m_{sq}^4} \right] ,
\]

with

\[
\tan(\beta) = v_2/v_1 ,
\]

\[
w = v_t^2 + v_2^2 \simeq (174 \text{ GeV})^2 ,
\]

\[
\Delta = \lambda A_\lambda < S > + \lambda \kappa < S >^2 - \frac{3m_t^2 A_t \lambda < S >}{16\pi^2 v_2^2} \ln \left( \frac{m_{sq}^2}{M_Z^2} \right) .
\]

Actually, neglecting the trilinear couplings \( A_t \) and \( A_\lambda \), in the leading log approximation and in the limit \( \tan(\beta) \to \infty \) the following analytic expression for the upper bound on the mass squared \( m_H^2 \) of the lightest Higgs scalar can be given \(^{23}\):

\[
m_H^2 \leq M_Z^2 \left[ 1 - \sin^2(2\beta) + \frac{2\lambda^2}{g_1^2 + g_2^2} \sin^2(2\beta) \right] + \frac{3}{4\pi^2 v_2^4 g_t^4} \ln \left( \frac{m_{sq}^2}{m_t^2} \right) .
\]

We have found numerically, that for non-vanishing \( A_t \) and \( A_\lambda \) and for arbitrary vacuum expectation value \( < S > \) and \( \kappa \) the bound \(^{19}\) is exceeded by at most 10 GeV provided \( A_t \) and \( A_\lambda \) are bounded by 1 TeV (in agreement with the observation made in \(^{11}\)).

From eqs. \(^{11}\) and the mass matrix \( M_{ij} \) (or the approximate result \(^{19}\)) it is evident that upper limits on the couplings \( g_t \) and \( \lambda \) turn into upper limits on \( m_t \) and \( m_H^2 \). Upper limits on \( g_t \) and \( \lambda \) can be obtained from the assumption that the running couplings develop no Landau singularities up to a certain scale \( \Lambda \) with, e.g., \( \Lambda \simeq 10^{16} \text{ GeV} \). These limits
have recently been studied in [24–28]. (According to [25] and [26] the limits of [24], which
have been obtained using analytic approximative solutions for the running couplings, are
somewhat too stringent.) From [25], e.g., one finds for \( g_t \) for general \( \kappa \) in eq. (10),
\[
g_t \leq 1.13 ,
\]  
whereas the upper bound on \( \lambda \) varies with \( g_t \). For \( g_t > 0.5 \) one finds (see also [21])
\[
\lambda \lesssim 0.87 .
\]  
The bound (20) translates into
\[
m_t \leq 195 \text{ GeV} .
\]  
A saturation of this bound implies actually a maximization of \( v_2 \); explicitly we have with
eq (11) and the bound (20) for fixed \( m_t \)
\[
v_2^2 \gtrsim \frac{m_t^2}{(1.13)^2} ,
\]  
or, with \( w = v_1^2 + v_2^2 \) kept fixed,
\[
\tan^2 \beta \gtrsim \frac{m_t^2}{(1.13)^2 w - m_t^2} .
\]  
In the evaluation of the upper bound on \( m_H \) according to eqs. (14) – (13) we will make
no further assumption on \( \beta \). Then one finds that, for \( m_t \gtrsim 130 \text{ GeV} \), the upper bound
on the lightest eigenvalue of \( M_{ij} \) is maximized by minimizing \( \tan^2 \beta \). (Note that the last
term on the right hand side of eq. (19) can be written as \( \frac{3m_t^4 (1 + \tan^2 \beta)}{4w^2w \tan^2 \beta} \ln \left( \frac{m_t^2}{m_t^2} \right) \).) Hence, for
\( m_t \gtrsim 130 \text{ GeV} \), we can fix \( \beta \) by saturating the bound (24). This expresses the fact that,
for increasing \( m_t, v_2 \) and hence \( \tan \beta \) have to increase in order not to violate the bound (20). Accordingly, whereas the contributions due to the radiative corrections increase with
\( m_t, \sin^2 2\beta \) decreases with \( m_t \). This implies that for large \( m_t \) (where \( \tan \beta \) has to be large)
the tree level contribution to \( m_H^2 \) proportional to \( \lambda^2 \) becomes negligible. Thus in this
region the upper bound on the lightest Higgs mass is the same as in the minimal model
including radiative corrections, which have to be computed respecting the bound (20). A
corresponding observation has also been made in [28], where models with additional singlets
and triplets have been considered (and bounds similar to ours have been obtained).

From a numerical analysis we find with the upper bounds of [25] for \( \lambda \) and (20) for \( g_t \)
that the upper limit on \( m_H \) varies between 140 GeV and 165 GeV \(^4\) for maximal \( A_t, A_\lambda \)
and \( m_{eq} \) of 1 TeV. The top mass, in turn, is bounded by 195 GeV (22). The combined

\(^4\) Here 10 GeV have been added in order to take care of the neglected effects discussed above.
limits on $m_H$ and $m_t$ surround the areas around the origin in Fig. 1 denoted by SUSY and SM+SUSY.

In the non–supersymmetric case we show the full numerical solution of the bound corresponding to (9) with $\Lambda = 1$ $TeV$ due to the unsolved hierarchy problem. The correspondingly allowed area in Fig. 1 is marked by SM or SM+SUSY. The two cases differ significantly and lead to areas in $m_H-m_t$ parameter space which are exclusively pointing into one of the two directions (these areas are labeled SM and SUSY, respectively). There are also areas, however, where both or neither of the scenarios are acceptable. These areas are denoted in Fig. 1 by SM+SUSY and NEITHER.

As already mentioned it is in principle conceivable that there exists a solution to the hierarchy problem without involving supersymmetry. In that case one could make $\Lambda$ within the Standard Model scenario very large, for example $10^{15}$ $GeV$. This would imply that the bounds for the Standard Model scenario would become significantly stronger. We have included in Fig. 1 these stronger Standard Model bounds as a weak solid line labelled $10^{15}$. Note that the area labeled NEITHER would then grow significantly and the area labeled SUSY exclusively would also grow, while the pure SM range as well as the SM+SUSY range would shrink. We have also included the experimental lower bounds on the Higgs mass [32] and the top mass [33] as dashed–dotted line.

The bounds of ref. [30] would be obtained if we would restrict our discussion to the minimal supersymmetric scenario, if we would ignore simultaneously the bound on $m_t$ from eq. (22), and if we took $\Lambda = 10^{15}$ $GeV$ for the Standard Model bounds without consideration of the hierarchy problem. We think, however, that it is better to take our more general scenario as the testing ground of supersymmetry, to implement the triviality constraint on $g_t$ as well, and that one should most likely take $\Lambda \simeq 1$ $TeV$ for the non–supersymmetric scenario. Consequently our bounds differ significantly from those of ref. [30].

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References

[1] I.V. Krive and A.D. Linde, Nucl. Phys. B117 (1976) 265; 
P.Q. Hung, Phys. Rev. Lett. 42 (1979) 873; 
H.D. Politzer and S. Wolfram, Phys. Lett. B82 (1979) 242; 
N.V. Krasnikov, Yad. Fiz. 28 (1978) 549; 
A.A. Anselm, JETP Lett. 29 (1979) 590.

[2] J. Glimm and A. Jaffe, Ann. Inst. H. Poincare 22 (1975) 97; 
J. Fröhlich, Nucl. Phys. B200 (1982) 281; 
M. Aizenmann, Phys. Rev. Lett. 47 (1981) 577; 
W.A. Bardeen and M. Moshe, Phys. Rev. D28 (1983) 1372.

[3] L.D. Landau and I. Pomeranchuk, Dokl. Akad. Nauk USSR 102 (1955) 489.

[4] N. Cabbibo, L. Maiani, G. Parisi and R. Petronzio, Nucl. Phys. B158 (1979) 296.

[5] M. Lindner, Z. Phys. C31 (1986) 295.

[6] M. Lindner, M. Sher and H. Zaglauer, Phys. Lett. B228 (1989) 139.

[7] M. Sher, Phys. Rep. 179 (1989) 273.

[8] U. Ellwanger, Fortschr. Phys. 36 (1988) 881.

[9] W.A. Bardeen, C.T. Hill and M. Lindner, Phys. Rev. D41 (1990) 1647.

[10] Y. Okada, M. Yamaguchi and T. Yanagida, Prog. Theor. Phys. Lett. 85 (1991) 1 and 
    Phys. Lett. B262 (1991) 54.

[11] J. Ellis, G. Ridolfi and F. Zwirner, Phys. Lett. B257 (1991) 83.

[12] H. Haber and R. Hempfling, Phys. Rev. Lett. 66 (1991) 1815.

[13] R. Barbieri, M. Frigeni and F. Caravaglios, Phys. Lett. B258 (1991) 167.

[14] J. Ellis, G. Ridolfi and F. Zwirner, Phys. Lett. B262 (1991) 477.

[15] A. Yamada, Phys. Lett. B263 (1991) 233.

[16] J. Espinosa and M. Quiros, Phys. Lett. B266 (1991) 389.

[17] P. Chankowski, S. Pokorski and J. Rosiek, Phys. Lett. B274 (1992) 191.

[18] A. Brignole, Phys. Lett. B281 (1992) 284.
[19] H. Nilles, M. Srednicki and D. Wyler, Phys. Lett. B120 (1983) 346.

[20] J.-P. Derendinger and C. Savoy, Nucl. Phys. B237 (1984) 307.

[21] J. Ellis, J. Gunion, H. Haber, L. Roszkowski and F. Zwirner, Phys. Rev. D39 (1989) 844.

[22] M. Drees, Int. Journ. Mod. Phys. A4 (1989) 3635.

[23] U. Ellwanger and M. Rausch de Traubenberg, Z. Phys. C53 (1992) 521.

[24] P. Binetruy and C. Savoy, Phys. Lett. B277 (1992) 453.

[25] J. Espinosa and M. Quiros, Phys. Lett. B279 (1992) 92.

[26] W. ter Veldhuis, Purdue preprint PURD-TH-92-11

[27] D. Comelli, INFN preprint (Nov. 1992).

[28] J. Espinosa and M. Quiros, preprint IEM-FT-60/92.

[29] R. Flores and M. Sher, Ann. Phys. 148 (1983) 95.

[30] N. Krasnikov and S. Pokorski, Phys. Lett. B288 (1992) 184.

[31] S. Bornholdt and C. Wetterich, Phys. Lett. B282 (1992) 399.

[32] L. Ronaldi, plenary talk at the 26th Conf. on High Energy Physics, Dallas, USA, 1992.

[33] CDF collaboration, F. Abe et al., Phys. Rev. D43 (1991) 2070.

**Figure Captions**

1. Combination of Higgs and top mass bounds of our general supersymmetric scenario (the fat solid line connecting \((m_H, m_t) = (113 \text{ GeV}, 0 \text{ GeV})\) with \((m_H, m_t) = (0 \text{ GeV}, 174 \text{ GeV})\)) with the Standard Model bounds for \(\Lambda = 1 \text{ TeV}\) (the fat solid line starting at \((m_H, m_t) = (0 \text{ GeV}, 85.6 \text{ GeV})\)). The resulting four areas are labeled with **NEITHER**, **SM**, **SUSY** and **SM+SUSY** respectively to indicate the allowed scenario(s). The weak solid line shows the stronger Standard Model bound for \(\Lambda = 10^{16} \text{ GeV}\). Experimental lower limits for the Higgs and top masses are shown as weak dashed–dotted lines.
Figure 1: Combination of Higgs and top mass bounds of our general supersymmetric scenario (the fat solid line connecting $(m_H, m_t) = (113 \text{ GeV}, 0 \text{ GeV})$ with $(m_H, m_t) = (0 \text{ GeV}, 174 \text{ GeV})$) with the Standard Model bounds for $\Lambda = 1 \text{ TeV}$ (the fat solid line starting at $(m_H, m_t) = (0 \text{ GeV}, 85.6 \text{ GeV})$). The resulting four areas are labeled with **NEITHER**, **SM**, **SUSY** and **SM+SUSY** respectively to indicate the allowed scenario(s). The weak solid line shows the stronger Standard Model bound for $\Lambda = 10^{15} \text{ GeV}$. Experimental lower limits for the Higgs and top masses are shown as weak dashed–dotted lines.