The upper bound on the lightest Higgs mass in the NMSSM revisited

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Abstract

We update the upper bound on the lightest CP even Higgs mass in the NMSSM, which is given as a function of $\tan\beta$ and $\lambda$. We include the available one and two loop corrections to the NMSSM Higgs masses, and constraints from the absence of Landau singularities below the GUT scale as well as from the stability of the NMSSM Higgs potential. For $m_{t_{\text{top}}}$ varying between 171.4 and 178 GeV, squark masses of 1 TeV and maximal mixing the upper bound is assumed near $\tan\beta \sim 2$ and varies between 139.9 and 141.4 GeV.

PAC numbers: 12.60.Jv, 14.80.Cp, 14.80.Ly

December 2006
LPT Orsay 06-77
LPTA Montpellier 06-62

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1 Introduction

Supersymmetric extensions of the standard model predict quite generally at least one relatively light Higgs boson. Hence, as soon as results from future collider experiments provide us with informations on the mass of at least one Higgs boson, we will be able to put constraints on possible supersymmetric extensions of the standard model.

To this end we need to know, as accurately as possible, how the Higgs boson masses depend on the nature of the supersymmetric extensions of the standard model, and on the parameters of these models. (We hope, of course, to get independent informations on these parameters from direct sparticle detections in the future.)

In the MSSM, corresponding calculations have been pushed to a fairly high accuracy, including many two-loop corrections. Recent reviews on the lightest Higgs boson mass in the MSSM can be found in refs. [1–3].

In the present paper we discuss the simplest version of the NMSSM [4] with a scale invariant superpotential

\[ W = \lambda S H_u H_d + \frac{\kappa}{3} S^3 + \ldots, \] (1.1)

which is the only supersymmetric extension of the Standard Model where the weak scale originates from the soft susy breaking scale only, i.e. where no supersymmetric dimensionful parameters as \( \mu \) are present in the superpotential.

It is well known [4] that the lightest Higgs boson in the NMSSM can be heavier than the one of the MSSM due to additional terms in the tree level Higgs potential proportional to \( \lambda^2 \); the additional contribution is

\[ \Delta m_h^2 = \frac{\lambda^2}{g^2} M_Z^2 \sin^2 2\beta. \] (1.2)

If one requires the absence of a Landau singularity for \( \lambda \) below the GUT scale, \( \lambda \) is bounded by \( \sim 0.7 \) from above [4], leading still to an upper bound on the mass of the lightest Higgs boson that is, however, larger than in the MSSM. Thus, future measurements of the Higgs boson mass could serve to distinguish these two models, provided that we know the difference between the upper bound on the mass of the lightest Higgs boson in the different models.

At present, the radiative corrections in the NMSSM have not been computed to quite the same accuracy as in the MSSM. Of course, radiative corrections in the NMSSM that are proportional to the quark/lepton Yukawa couplings and the gauge couplings only are the same as in the MSSM, but there are many additional contributions involving the new Yukawa couplings \( \lambda \) and \( \kappa \) in the superpotential in eq. (1.1), and the associated soft trilinear couplings \( A_\lambda \) and \( A_\kappa \).
The one loop corrections in the NMSSM induced by t and b quark/squark loops have been computed already some time ago [5], and the dominant two loop corrections ($\sim h_t^6$ and $\sim h_4^2\alpha_s$), that are the same as in the MSSM, have been included in an analysis of the NMSSM Higgs sector in ref. [6, 7].

The leading logarithmic one loop corrections to the lightest Higgs mass in the NMSSM proportional to the electroweak gauge couplings $g$ or NMSSM specific Yukawa couplings $\lambda$, $\kappa$ ($\sim g^4, g^2\lambda^2, g^2\kappa^2, \lambda^4, \kappa^4$) have been computed only recently [8]. They are included in the latest version of the code NMHDECAY [9–11], where the NMSSM Higgs masses, couplings and branching ratios are computed as functions of the parameters in the Lagrangian of the model.

This code checks also the absence of a Landau singularity for $\lambda$ below the GUT scale using the two loop renormalization group equations, and susy threshold effects around the susy scale. This procedure is numerically relevant, since $\Delta m_h^2$ depends on $\lambda$, and the upper bound on $\lambda$ depends on $\tan\beta$ (via the top Yukawa coupling $h_t$) and $\kappa$.

It is the purpose of the present paper to review the upper bound on the lightest Higgs boson mass in the NMSSM, using the up-to-date knowledge of the corresponding radiative corrections.

Instead of investigating far-fetched regions in parameter space that serve to obtain very conservative bounds, we proceed as follows: For the soft terms that are relevant for the sparticle spectrum, we chose universal squark and slepton masses of 1 TeV, and trilinear couplings of 2.5 TeV (that practically maximize the one loop radiative corrections to the Higgs boson mass). For the gaugino masses we take $M_1 = 150$ GeV, $M_2 = 300$ GeV and $M_3 = 1$ TeV in rough agreement with universal gaugino masses at the GUT scale. For the top quark pole mass we present results both for $m_{top} = 171.4$ GeV (the latest central value obtained by the Tevatron Electroweak Working Group [12]) and a very conservative upper limit of $m_{top} = 178$ GeV. The NMSSM specific Yukawa couplings and trilinear soft terms $\lambda$, $\kappa$, $A_\lambda$ and $A_\kappa$ as well as the effective $\mu$ parameter ($\mu = \lambda s$ in the NMSSM) are chosen such that the lightest Higgs boson mass is maximized, without violating constraints from perturbativity of the Yukawa couplings at the GUT scale, nor phenomenological constraints on CP odd or charged Higgs masses and couplings. To this end a numerical analysis is required, that is performed using the updated version of NMHDECAY [11]. The upper bound on the lightest Higgs mass is then given as a function of $\tan\beta$ and $\lambda$.

For the same choice of the above soft terms, we present the upper bound on the lightest Higgs mass in the MSSM limit $\lambda \rightarrow 0$ as obtained with NMHDECAY. This result can be compared to values obtained from analytical or numerical analyses in the MSSM, that include
radiative corrections that are still absent in NMHDECAY: these are notably electroweak one-loop corrections $\sim g^4$ beyond the LLA, and non-dominant two-loop corrections (involving less than two powers of large logarithms) $\sim h_t^6$ and $\sim h_t^4 \alpha_s$ beyond the ones that follow from an RG improvement of the one loop corrections [13] (which are included).

For soft terms as above, $m_{t_{\text{top}}} = 178$ GeV, $M_A = \mu = 1$ TeV and $\tan\beta = 10$ SuSpect gives $m_h \sim 128.5$ GeV (taken from ref. [2]), FeynHiggs $m_h \sim 134$ GeV (taken from ref. [3]), and NMHDECAY $m_h \sim 128.6$ GeV. This allows to estimate the uncertainties on $m_h$ due to the radiative corrections not included in NMHDECAY, following the discussions in [2, 3] that we will not repeat here.

The striking effect in the NMSSM is that the maximal value of $m_h$ is not assumed for large $\tan\beta$ as in the MSSM, but at low $\tan\beta \sim 2$ due to the tree level term noted above. There we obtain $m_h \sim 139.9$ GeV for $m_{t_{\text{top}}} = 171.4$ GeV, and $m_h \sim 141.4$ GeV for $m_{t_{\text{top}}} = 178$ GeV (for the same other parameters as above). For larger values of $\tan\beta$ the upper bound on $m_h$ decreases in the NMSSM. For $\tan\beta \gtrsim 10$ it hardly exceeds the MSSM value given above, since the effect of the tree level term becomes small. For small $\tan\beta \lesssim 2$ the absence of a Landau singularity below $M_{\text{GUT}}$ restricts $\lambda$ more strongly from above, due to the large top Yukawa coupling $h_t$. This implies that present lower limits on $m_h$ from LEP still lead to a lower bound on $\tan\beta$ of $\sim 1.3$ in the NMSSM.

It must be noted that in particular regions of the parameter space of the NMSSM the upper bound on $m_h$ discussed here can be misleading:

In principle, a singlet-like CP even Higgs boson can be lighter than the lightest doublet-like CP even Higgs boson (with non-vanishing couplings to the $Z$ boson) in the NMSSM. Strictly speaking, the upper bound on the lightest CP even Higgs boson discussed here is then still valid.

However, a singlet-like CP even Higgs boson would have been practically undetectable at LEP due to its vanishing coupling to the $Z$ boson. Fortunately, if the lightest CP even Higgs boson is a pure singlet in the NMSSM, the upper bound on $m_h$ discussed here applies then to the lightest doublet-like CP even Higgs boson. On the other hand, if the lightest CP even Higgs boson is only approximately a singlet, the lightest doublet-like CP even Higgs boson can be heavier than the upper bound on $m_h$ discussed here.

A similar reasoning applies to the situation where the doublet-like CP even (SM like) Higgs boson decays into singlet like (mostly two CP odd) scalars [14]. Then, the detection of the SM like Higgs boson can be very challenging, even if its mass satisfies the upper bounds discussed here. Hence, although the upper bound on $m_h$ presented here is always valid, it may refer to a state that is difficult to detect.
2 The upper bound on the lightest Higgs Boson mass in the NMSSM

In order to find the regions in the parameter space of the NMSSM that maximize the upper bound on the lightest CP even Higgs Boson mass, it is helpful to take a look at the CP even Higgs mass matrix at tree level. In the basis \((H_u, H_d, S)\) and using the minimization equations in order to eliminate the soft masses squared, it reads:

\[
\mathcal{M}_S^2 = \begin{pmatrix}
    g^2 h_u^2 + \mu \frac{h_d}{h_u} (A_\lambda + \nu) & (2\lambda^2 - g^2) h_u h_d - \mu (A_\lambda + \nu) & 2\lambda h_u \mu - \lambda h_d (A_\lambda + 2\nu) \\
    (2\lambda^2 - g^2) h_u h_d & g^2 h_d^2 + \mu \frac{h_u}{h_d} (A_\lambda + \nu) & 2\lambda h_d \mu - \lambda h_u (A_\lambda + 2\nu) \\
    \lambda^2 A_\lambda \frac{h_u h_d}{\mu} + \nu (A_\kappa + 4\nu) & 2\lambda h_d \mu - \lambda h_u (A_\lambda + 2\nu) & 0
\end{pmatrix}
\]  

(2.1)

where \(\nu = \kappa s\). To a good approximation, the \(2 \times 2\) doublet subsector is diagonalized by the angle \(\beta\) which gives the desired light eigenstate \(h\) and a heavy eigenstate \(H\) with a mass \(m_H \sim m_A\) close to the MSSM-like CP odd state (the larger \(m_A\), the better this approximation). In the NMSSM, one can define \(m_A^2\) as the diagonal doublet term in the CP odd \(2 \times 2\) mass matrix after the Goldstone mode has been dropped. At tree level, it has the same expression as in the MSSM:

\[
m_A^2 = \frac{2 \mu B}{\sin 2\beta} , \quad \text{with} \quad B = A_\lambda + \nu .
\]  

(2.2)

In the CP even sector of the NMSSM this is not the end of the story, however: the light eigenstate \(h\) of the \(2 \times 2\) doublet subsector still mixes with the singlet state \(S\), which is heavier than \(h\) by assumption. In order to maximize \(m_h\), this mixing has to vanish:

\[
\lambda [2\mu - (A_\lambda + 2\nu) \sin 2\beta] \sim 0 .
\]  

(2.3)

This requires either \(\lambda \to 0\) (which minimizes the NMSSM specific tree level contribution to \(m_h\) of eq. (1.2)) or

\[
A_\lambda \simeq \frac{2\mu}{\sin 2\beta} - 2\nu .
\]  

(2.4)

On \(A_\kappa\) we get the following constraints: \(\mathcal{M}_{S,33}^2\) in eq. (2.1) must at least be positive, which requires essentially (since the first term is typically relatively small) \(A_\kappa \nu \gtrsim -4\nu^2\). The CP odd mass matrix element in the singlet sector, given by

\[
\mathcal{M}_{P,33}^2 = 4\lambda \kappa h_u h_d + \lambda^2 A_\lambda \frac{h_u h_d}{\mu} - 3\nu A_\kappa ,
\]  

(2.5)
must also be positive. Typically the last term in (2.5) dominates, hence we get an allowed window

\[-4\nu^2 \lesssim A_\kappa \nu \lesssim 0.\]  

(2.6)

Next, in order to maximize the NMSSM specific tree level contribution to \(m_h\) of eq. (1.2), \(\lambda\) has to be as large as possible; we require, however, the absence of a Landau singularity for all Yukawa couplings \(\lambda, \kappa, h_t\) below the GUT scale, which leads to the following constraints:

First, given the corresponding RG equations [4], this implies small values for \(\kappa\). The limit \(\kappa \to 0\) while \(\lambda\) remains finite is disallowed, however, both from the stability of the potential and the fact that the allowed window of eq. (2.6) vanishes in this limit. (Of course, stability of the potential and positivity of all masses squared are related issues.)

Second, for small \(\tan\beta\) the top quark Yukawa coupling becomes large, and can run into a Landau singularity below \(M_{GUT}\), or induce a Landau singularity below \(M_{GUT}\) for \(\lambda\). The value of \(\tan\beta\) that allows for maximal values of \(\lambda\) (and maximizes the tree level contribution to \(m_h\) of eq (1.2)) is around 2.

In this region of \(\tan\beta\), a larger value for the top quark pole mass does hardly increase the upper bound on \(m_h\): at fixed \(\tan\beta\), larger \(m_{top}\) implies a larger top Yukawa coupling \(h_t\), which implies a somewhat lower allowed value for \(\lambda\). Consequently a variation of the top quark pole mass between 171.4 and 178 GeV (which increases \(m_h\) by \(\sim 4.8\) GeV for \(\tan\beta \sim 10\)), increases the maximal allowed value for \(m_h\) in the NMSSM by only \(\sim 1.5\) GeV for \(\tan\beta \sim 2\).

For large values of \(\tan\beta\), it is obvious from eq. (2.2) that \(m_A\) tends to be very large, unless \(B\) is small. (\(\mu\) cannot be smaller than \(\sim 100\) GeV due to the lower bound on chargino masses from LEP). Very large values of \(m_A\) are unnatural, since they require supersymmetry breaking Higgs masses of the same order of magnitude, which aggravate the fine tuning problem – a situation which we want to avoid. In the MSSM, one can always chose \(B\) small enough to keep \(m_A\) reasonable even at large \(\tan\beta\). In the NMSSM this is also possible, provided that \(\nu \simeq -A_\lambda\). However, one has also to minimize the doublet-singlet mixing of eq. (2.3) in order to to maximize \(m_h\). If \(\lambda\) is not very small, eq. (2.4) together with \(\nu \simeq -A_\lambda\) implies \(\nu \gg \mu\), which is equivalent to \(\kappa \gg \lambda\). Large values of \(\kappa\) leading to a Landau singularity below the GUT scale, this is excluded. Thus, the only way of minimizing the doublet-singlet mixing while keeping \(m_A\) constant at large \(\tan\beta\) is to assume \(\lambda \to 0\), which means that the bound on \(m_h\) is the same as in the MSSM.

(In general, for large values of \(\tan\beta\) the LEP constraints on \(m_h\) imply either that \(m_A\) and \(|A_\lambda|\) assume very large values \(\gtrsim 1\) TeV, or \(\lambda \lesssim 0.2\).)
All these considerations make it clear that a realistic upper limit on \(m_h\) in the NMSSM requires numerical methods; analytic approaches can be misleading (and can allow for larger values of \(m_h\)).

Our results below are obtained with \texttt{NMHDECAY} \cite{11}. The precision of the included radiative corrections to the lightest CP even Higgs mass has already been discussed in the introduction and is given in \cite{9} and \cite{10}.

As discussed in the introduction, we take universal squark and slepton masses of 1 TeV, and trilinear squark/slepton couplings of 2.5 TeV (near maximal mixing). For the gaugino masses we take \(M_1 = 150\) GeV, \(M_2 = 300\) GeV and \(M_3 = 1\) TeV. We scan over the NMSSM specific Yukawa couplings and trilinear soft terms \(\lambda, \kappa, A_\lambda, A_\kappa\) as well as the effective \(\mu\) parameter, and we obtain the regions in the NMSSM parameter space that maximize \(m_h\) in agreement with the considerations above.

In fig. 1 we show our results for the the upper bound on \(m_h\) for \(1 < \tan \beta < 10\). The thick full line corresponds to \(m_{\text{top}} = 178\) GeV, the thin full line to \(m_{\text{top}} = 171.4\) GeV, both without imposing constraints on \(m_A\).

With the above soft terms, the upper bound on \(m_h\) in the NMSSM is 141.4 GeV for \(m_{\text{top}} = 178\) GeV. It is reached for \(\tan \beta \sim 2.2, \lambda \sim .677, \kappa \sim .068, \mu \sim 545\) GeV, \(A_\lambda \sim 1365\) GeV, and \(A_\kappa \sim 10\) GeV (strictly speaking a certain range of values for \(\kappa, \mu, A_\lambda\) and \(A_\kappa\) gives the same result for \(m_h\) for these values of \(\tan \beta\) and \(\lambda\)). For \(m_{\text{top}} = 171.4\) GeV, the upper bound on \(m_h\) is 139.9 GeV and is obtained for \(\tan \beta \sim 2, \lambda \sim .703, \kappa \sim .049, \mu \sim 534\) GeV, \(A_\lambda \sim 1287\) GeV and \(A_\kappa \sim 10\) GeV.

For \(\tan \beta = 10\) we get 133.6 GeV in the NMSSM for \(m_{\text{top}} = 178\) GeV (resp. 128.8 GeV for \(m_{\text{top}} = 171.4\) GeV), which remains nearly constant for larger values of \(\tan \beta\) (a slight increase of the contributions from the radiative corrections is compensated by a slight decrease of the tree level term of eq. (1.2)).

In the same fig. 1, we show the upper bound on \(m_A\) in the MSSM limit \(\lambda \to 0\) as obtained with \texttt{NMHDECAY} as a thick dashed line for \(m_{\text{top}} = 178\) GeV, and as a thin dashed line for \(m_{\text{top}} = 171.4\) GeV (taking \(m_A = 1\) TeV). In this limit, the upper bound on \(m_h\) reaches 129.7 GeV for \(m_{\text{top}} = 178\) GeV (resp. 124.4 GeV for \(m_{\text{top}} = 171.4\) GeV) at \(\tan \beta = 10\), and increases by another 1 GeV for very large \(\tan \beta = 50\).

As noted above, large values of \(\tan \beta\) imply large values for \(m_A\) in the NMSSM, if \(\lambda\) is kept fixed. Indeed, along the full lines of fig. 1 the value of \(m_A\) increases with \(\tan \beta\) up to several TeV. The consequence of fixing \(m_A \leq 1\) TeV is that the maximally allowed value of \(\lambda\) decreases with \(\tan \beta\). The corresponding effect on the upper bound of \(m_h\) is shown as a thick dotted line for \(m_{\text{top}} = 178\) GeV, and a thin dotted line for \(m_{\text{top}} = 171.4\) GeV in
Figure 1: Upper bound on the lightest Higgs mass in the NMSSM for $m_{\text{top}} = 178$ GeV (thick full line: $m_A$ arbitrary, thick dotted line: $m_A = 1$ TeV) and $m_{\text{top}} = 171.4$ GeV (thin full line: $m_A$ arbitrary, thick dotted line: $m_A = 1$ TeV) and in the MSSM (with $m_A = 1$ TeV) for $m_{\text{top}} = 178$ GeV (thick dashed line) and $m_{\text{top}} = 171.4$ GeV (thin dashed line) as obtained with NMHDECAY as a function of $\tan\beta$. Squark and gluino masses are 1 TeV and $A_{\text{top}} = 2.5$ TeV.

Let us compare this bound on $m_h$ to earlier work: it is about 6 GeV larger than the one obtained from fig. 4 in ref. [7] (for the corresponding values for $m_{\text{top}}$). Also the value of $\tan\beta$, where this bound is reached, is now smaller ($\sim 2$ compared to $\sim 3$ in ref. [7]).
Figure 2: Upper bound on the lightest Higgs mass in the NMSSM for $m_{\text{top}} = 178$ GeV, $\tan \beta = 2.2$, electroweak/Yukawa corrections included (thick full line) and omitted (thick dotted line), and $m_{\text{top}} = 171.4$ GeV, $\tan \beta = 2$, electroweak/Yukawa corrections included (thin full line) and omitted (thin dotted line). Squark and gluino masses and $A_{\text{top}}$ are as in fig. 1.

Differences are due to the improved treatment of radiative corrections in NMHDECAY which concerns both the two loop corrections $\sim h_t^6$ and $h_t^4 \alpha_s$ (which are now RG-improved), and the inclusion of one loop corrections (in the LLA, keeping terms $\sim \ln(M_{\text{Susy}}^2/M_Z^2)$) proportional to the electroweak gauge couplings and NMSSM specific Yukawa couplings $\lambda$ and $\kappa$. The effect of the first improvement is a considerable increase in $m_h$, whereas the effect of the electroweak/Yukawa corrections is a slight decrease of $m_h$ by up to $\sim 2$ GeV.

In order to clarify the latter effect and, simultaneously, the general effect of the NMSSM specific Yukawa couplings at low $\tan \beta$, we show in fig. 2 the upper bound on $m_h$ as a function of $\lambda$ at fixed $\tan \beta$. Here the thick full line corresponds to $m_{\text{top}} = 178$ GeV, $\tan \beta = 2.2$ and electroweak/Yukawa corrections included, whereas the thick dotted line would be the result with these corrections omitted. The thin full line corresponds to $m_{\text{top}} = 171.4$ GeV, $\tan \beta = 2$ and electroweak/Yukawa corrections included, whereas the thin dotted line would be the
result without these corrections. One sees the decrease in $m_h$ due to the electroweak/Yukawa corrections, which increases the lower bound on $\lambda$ for small values of $\tan \beta$ and $m_{top}$ due to the LEP bound on $m_h$.

As final remark we repeat, as noted at the end of the introduction, that the mass of the lightest detectable Higgs boson could be larger in the NMSSM than the upper bounds given here; in order to interpret future data in the context of the NMSSM, constraints (or positive results) must be available in the plane Higgs mass versus Higgs couplings in order to be sensitive to a possible singlet/doublet mixing.

**Acknowledgement**

We like to thank P. Slavich for helpful discussions, notably for pointing out a mistake in the first version of this paper. We acknowledge support by the ANR grant PHYS@COL&COS.
References

[1] M. Carena and H. E. Haber, Prog. Part. Nucl. Phys. 50 (2003) 63 [arXiv:hep-ph/0208209].

[2] B. C. Allanach, A. Djouadi, J. L. Kneur, W. Porod and P. Slavich, JHEP 0409 (2004) 044 [arXiv:hep-ph/0406166].

[3] S. Heinemeyer, W. Hollik and G. Weiglein, Phys. Rept. 425 (2006) 265 [arXiv:hep-ph/0412214].

[4] H. P. Nilles, M. Srednicki and D. Wyler, Phys. Lett. B 120 (1983) 346;
J. M. Frere, D. R. T. Jones and S. Raby, Nucl. Phys. B 222 (1983) 11;
J. P. Derendinger and C. A. Savoy, Nucl. Phys. B 237 (1984) 307;
J. R. Ellis, J. F. Gunion, H. E. Haber, L. Roszkowski and F. Zwirner, Phys. Rev. D 39 (1989) 844;
M. Drees, Int. J. Mod. Phys. A 4 (1989) 3635;
U. Ellwanger, M. Rausch de Traubenberg and C. A. Savoy, Phys. Lett. B 315 (1993) 331 [hep-ph/9307322] and Nucl. Phys. B 492 (1997) 21 [hep-ph/9611251];
S. F. King and P. L. White, Phys. Rev. D 52 (1995) 4183 [hep-ph/9505326];
F. Franke and H. Fraas, Int. J. Mod. Phys. A 12 (1997) 479 [hep-ph/9512366].

[5] U. Ellwanger, Phys. Lett. B 303 (1993) 271 [arXiv:hep-ph/9302224];
T. Elliott, S.F. King, P.L. White, Phys. Lett. B 305 (1993) 71 [hep-ph/9302202],
Phys. Lett. B 314 (1993) 56 [hep-ph/9305282] and Phys. Rev. D 49 (1994) 2435 [hep-ph/9308309];
P. Pandita, Phys. Lett. B 318 (1993) 338 and Z. Phys. C 59 (1993) 575;
S. Ham, S. Oh, B. Kim, J. Phys. G 22 (1996) 1575 [hep-ph/9604243].

[6] G. K. Yeghiyan, “Upper bound on the lightest Higgs mass in supersymmetric theories,”
arXiv:hep-ph/9904488;

[7] U. Ellwanger and C. Hugonie, Eur. Phys. J. C 25 (2002) 297 [arXiv:hep-ph/9909260].

[8] U. Ellwanger and C. Hugonie, Phys. Lett. B 623 (2005) 93 [arXiv:hep-ph/0504269].
[9] U. Ellwanger, J. F. Gunion and C. Hugonie, JHEP 0502 (2005) 066 [arXiv:hep-ph/0406215].

[10] U. Ellwanger and C. Hugonie, Comput. Phys. Commun. 175 (2006) 290 [arXiv:hep-ph/0508022].

[11] U. Ellwanger and C. Hugonie, “NMSPEC: A Fortran code for the sparticle and Higgs masses in the NMSSM with GUT scale boundary conditions”, [arXiv:hep-ph/0612134].

[12] E. Brubaker et al. [Tevatron Electroweak Working Group], “Combination of CDF and D0 results on the mass of the top quark,” arXiv:hep-ex/0608032.

[13] R. Barbieri, M. Frigeni and F. Caravaglios, Phys. Lett. B 258 (1991) 167;
    J. R. Espinosa and M. Quiros, Phys. Lett. B 266 (1991) 389;
    H. E. Haber and R. Hempfling, Phys. Rev. D 48 (1993) 4280 [arXiv:hep-ph/9307201];
    J. Kodaira, Y. Yasui and K. Sasaki, Phys. Rev. D 50 (1994) 7035 [arXiv:hep-ph/9311366];
    R. Hempfling and A. H. Hoang, Phys. Lett. B 331 (1994) 99 [arXiv:hep-ph/9401219];
    J. A. Casas, J. R. Espinosa, M. Quiros and A. Riotto, Nucl. Phys. B 436 (1995) 3
    [Erratum-ibid. B 439 (1995) 466] [arXiv:hep-ph/9407389];
    M. Carena, J. R. Espinosa, M. Quiros and C. E. M. Wagner, Phys. Lett. B 355 (1995)
    209 [arXiv:hep-ph/9504316];
    M. Carena, M. Quiros and C. E. M. Wagner, Nucl. Phys. B 461 (1996) 407 [arXiv:hep-ph/9508343];
    H. E. Haber, R. Hempfling and A. H. Hoang, Z. Phys. C 75 (1997) 539 [arXiv:hep-ph/9609331];
    M. Carena, H. E. Haber, S. Heinemeyer, W. Hollik, C. E. M. Wagner and G. Weiglein,
    Nucl. Phys. B 580 (2000) 29 [arXiv:hep-ph/0001002].

[14] B. A. Dobrescu, G. Landsberg and K. T. Matchev, Phys. Rev. D 63 (2001) 075003
    [arXiv:hep-ph/0005308];
    B. A. Dobrescu and K. T. Matchev, JHEP 0009 (2000) 031 [arXiv:hep-ph/0008192];
    R. Dermisek and J. F. Gunion, Phys. Rev. Lett. 95 (2005) 041801 [arXiv:hep-ph/0502105],
    Phys. Rev. D 73 (2006) 111701 [arXiv:hep-ph/0510322] and “The NMSSM
close to the R-symmetry limit and naturalness in $h \rightarrow aa$ decays for $m(a) < 2m(b)$,” arXiv:hep-ph/0611142;
U. Ellwanger, J. F. Gunion and C. Hugonie, JHEP 0507 (2005) 041 [arXiv:hep-ph/0503203];
S. Moretti, S. Munir and P. Poulose, Phys. Lett. B 644 (2007) 241 [arXiv:hep-ph/0608233];
S. Chang, P. J. Fox and N. Weiner, JHEP 0608 (2006) 068 and “Visible cascade Higgs decays to four photons at hadron colliders,” arXiv:hep-ph/0608310;
T. Stelzer, S. Wiesenfeldt and S. Willenbrock, “Higgs at the Tevatron in Extended Supersymmetric Models,” arXiv:hep-ph/0611242;
K. Cheung, J. Song and Q. S. Yan, “Search for the Higgs boson into two pseudoscalar bosons at the LHC,” arXiv:hep-ph/0703149.