TWO-LOOP CORRECTIONS IN THE HIGGS SECTOR OF THE MSSM*

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(Received October 11, 2004)

We present a computation of the leading two-loop corrections to the MSSM Higgs boson masses and electroweak symmetry breaking conditions. The computation is performed in the effective potential approach and includes corrections controlled by the third-family Yukawa couplings and the strong gauge coupling. We discuss a renormalisation scheme that avoids unphysically large threshold effects associated with the bottom Yukawa couplings. We also discuss the implementation of our corrections in computer programs that compute the MSSM mass spectrum from a set of unified high-energy boundary conditions.

PACS numbers: 12.60.Jv, 14.80.Cp

1. Introduction

A crucial prediction of the Minimal Supersymmetric Standard Model (MSSM) [1] is the existence of at least one light Higgs boson [2] which, at the tree level, is bound to be lighter than the Z boson. If this upper bound was not significantly raised by radiative corrections, the failure of detecting this Higgs boson at LEP would have ruled out the MSSM as a viable theory for physics at the weak scale. However, it was first realised in Ref. [3] that the inclusion of the one-loop $\mathcal{O}(\alpha_t)$ corrections, which rise quartically with the mass of the top quark and logarithmically with the mass of its scalar superpartner, may push the lighter Higgs boson mass well above the tree-level bound.

In the subsequent years, an impressive theoretical effort has been devoted to the precise determination of the Higgs boson masses in the MSSM. A first step was to provide the full one-loop computation, performed in Refs. [4, 5].

* Presented at the final meeting of the European Network "Physics at Colliders", Montpellier, France, September 26–27, 2004.

A second step was the addition of the dominant two-loop corrections which involve the strongest couplings of the theory: the QCD coupling constant and the Yukawa couplings of the third-generation fermions (although the masses of the bottom quark and the τ lepton are relatively tiny compared to the top quark mass, the b and τ Yukawa couplings can be strongly enhanced for large values of $\tan\beta$, the ratio of the two Higgs vacuum expectation values). The leading logarithmic effects at two loops have been included via appropriate RGEs [6, 7], and the genuine two-loop corrections of $\mathcal{O}(\alpha_t\alpha_s)$ [8–12] have been evaluated in the limit of zero external momentum. Subsequently, the two-loop Yukawa corrections of $\mathcal{O}(\alpha_t^2)$ [8, 11, 13], $\mathcal{O}(\alpha_b\alpha_s)$ [14], $\mathcal{O}(\alpha_t\alpha_b + \alpha_b^2)$ [15], $\mathcal{O}(\alpha_\tau^2)$ and $\mathcal{O}(\alpha_b\alpha_\tau)$ [16] have been evaluated in the limit of zero external momentum. The tadpole corrections needed to minimise the effective scalar potential, $V_{\rm eff}$, have also been calculated at the one-loop [5, 17] and two-loop [15, 16, 18] levels for the strong coupling and the third-generation Yukawa couplings. Finally, the full two-loop corrections to the MSSM effective potential have been calculated [19], together with a first study of the two-loop corrections to the lighter Higgs boson mass controlled by the electroweak gauge couplings [20] and of the leading momentum-dependent two-loop corrections [21].

The calculation of the radiative corrections to physical observables requires the choice of a renormalisation scheme for the input parameters. For example, the corrections can be expressed in terms of "On-Shell" (OS) parameters, such as pole particle masses and suitably defined mixing angles. This is straightforward when only the corrections involving the top Yukawa couplings are taken into account; however, it is well known that in the MSSM the relation between the physical bottom mass and the corresponding Yukawa coupling can receive large, $\tan\beta$-enhanced threshold corrections [22]. If the physical bottom mass is used as input parameter in the one-loop part of the computation, potentially large $\tan\beta$-enhanced corrections appear at two loops. To address this problem, a set of renormalisation prescriptions for the parameters in the bottom/sbottom sector that avoid the occurrence of unphysically large threshold effects at two loops was proposed in Refs. [14, 15] for the $\mathcal{O}(\alpha_b\alpha_s)$ and $\mathcal{O}(\alpha_t\alpha_b + \alpha_b^2)$ parts of the corrections. This OS renormalisation scheme has been adopted for the computation of the MSSM Higgs boson masses by the computer code `FeynHiggs` [23], based on the two-loop results of Refs. [9] and subsequently expanded to include the results of Refs. [13–15] (see Ref. [24] for a discussion).

In alternative, the MSSM input parameters can be expressed in a minimal renormalisation scheme such as modified dimensional reduction scheme, or $\overline{\rm DR}$. This way of presenting the results is convenient for analyzing models that predict, via the renormalisation group equations (RGE), the low-energy $\overline{\rm DR}$ values of the MSSM parameters in terms of a set of unified boundary conditions assigned at some scale $M_{\rm GUT}$ much larger than the weak scale. This is for instance the case of the models of gravity mediated (mSUGRA), gauge mediated (GMSB) or anomaly mediated (AMSB) SUSY breaking. Several computer codes have been developed in the past years to provide reliable determinations of the supersymmetric mass spectra in models with high-energy boundary conditions, including the calculations of the various radiative corrections. In particular, the publicly available codes `SoftSusy` [25], `SuSpect` [26] and `SPheno` [27] include a purely $\overline{\rm DR}$ calculation of the neutral MSSM Higgs boson masses, based on the one-loop results of Ref. [5] and the two-loop results of Refs. [12, 13, 15, 18] (for a detailed discussion see Ref. [16]).

This talk has two purposes: the first is to review the results of Refs. [14, 15], in which the leading two-loop corrections to the Higgs boson masses controlled by the bottom Yukawa coupling were computed and a suitable renormalisation scheme avoiding the appearence of unphysically large $\tan\beta$-enhanced terms was proposed. The second purpose is to review the results of Ref. [16], where the implementation of the two-loop corrections to the MSSM Higgs boson masses in the three public codes `SoftSusy`, `SuSpect` and `SPheno` was discussed.

2. On-shell renormalisation scheme for large $\tan\beta$

When they do not refer to boundary conditions at high scales, general analyses of the MSSM are usually performed in terms of parameters that allow for a direct physical interpretation, such as pole masses and appropriately defined mixing angles in the squark sector. It is rather easy to devise an OS renormalisation scheme for the parameters in the top/stop sector: we can use the OS prescription for the top and stop masses and the stop mixing angle, treat the trilinear stop interaction term A_t as a derived quantity and retain a $\overline{\rm DR}$ definition for the superpotential Higgs mass parameter μ and for $\tan\beta$ [13]. Instead, some additional care is required in the choice of an OS scheme for the parameters in the bottom/sbottom sector, due to the potentially large one-loop threshold corrections [22], proportional to $\tan\beta$, that contribute to the pole bottom mass. For example, a definition of A_b in terms of the OS bottom and sbottom masses and sbottom mixing angle, similar to the definition of A_t, would produce a counterterm δA_b proportional to $\tan^2\beta$ [28]. When $\tan\beta$ is large, this would induce very large corrections to the Higgs boson masses at two loops, questioning the validity of the perturbative expansion.

To overcome this problem, we adopt a set of renormalisation prescriptions for the parameters in the the bottom/sbottom sector that avoid the occurrence of unphysically large threshold effects and at the same time enforce other desirable properties such as the decoupling of heavy particles, the infrared finiteness and gauge-independence. Combining these prescriptions with the usual prescriptions for the top/stop parameters [13], we obtain a convenient OS renormalisation scheme for the $\mathcal{O}(\alpha_b\alpha_s)$ and $\mathcal{O}(\alpha_t\alpha_b+\alpha_b^2)$ parts of the corrections to the Higgs boson masses. Since the corrections controlled by the bottom Yukawa coupling can be sizeable only for large values of $\tan\beta$, we work directly in the physically relevant limit of $\tan\beta\to\infty$, *i.e.* $v_1\to 0\,,\ v_2\to v$.

The OS definitions for the squark masses and mixing angles, the top quark mass and the electroweak parameter $v\equiv(\sqrt{2}\,G_\mu)^{-1/2}$ fix the finite parts of the corresponding counterterms to:

$$\begin{aligned}\delta m^2_{\tilde q_i} &= \Pi^{\tilde q}_{ii}(m^2_{\tilde q_i})\,, \qquad \delta\theta_{\tilde q} = \frac{1}{2}\,\frac{\Pi^{\tilde q}_{12}(m^2_{\tilde q_1})+\Pi^{\tilde q}_{12}(m^2_{\tilde q_2})}{m^2_{\tilde q_1}-m^2_{\tilde q_2}}\,,\\ \delta m_t &= \Sigma_t(m_t)\,, \qquad\quad \delta v = \frac{v}{2}\,\frac{\Pi^T_{WW}(0)}{m^2_W}\,,\end{aligned}\tag{1}$$

where $\tilde q=(\tilde t\,,\tilde b)$, while $\Pi^{\tilde q}_{ij}(p^2)$, $\Sigma_t(p)$ and $\Pi^T_{WW}(p^2)$ denote the real and finite parts of the self-energies of squarks, top quark and W boson, respectively. Following Ref. [13], we further treat μ as a $\overline{\rm DR}$ parameter computed at a reference scale $Q_0=175$ GeV, and h_t and A_t as derived quantities that can be computed by means of the tree-level formulae for m_t and $s_{2\theta_t}$, respectively. In principle, we still have to define m_b, h_b and A_b. However, in the large $\tan\beta$ limit, the bottom mass is just zero, and the sbottom mixing angle becomes

$$s_{2\theta_b} = -\frac{\sqrt{2}\,h_b\,\mu\,v}{m^2_{\tilde b_1}-m^2_{\tilde b_2}}\,,\tag{2}$$

which is independent of m_b and A_b. We can thus treat h_b as a quantity derived from the sbottom mixing, and use Eqs. (1) and (2) to obtain a prescription for δh_b:

$$\delta h_b = h_b\left(\frac{\delta m^2_{\tilde b_1}-\delta m^2_{\tilde b_2}}{m^2_{\tilde b_1}-m^2_{\tilde b_2}}+\frac{\delta s_{2\theta_b}}{s_{2\theta_b}}-\frac{\delta v}{v}\right)\,.\tag{3}$$

Concerning the definition of A_b, we observe that the Yukawa coupling h_b multiplying A_b can be absorbed in a redefinition of the trilinear soft-breaking term, $\widetilde{A}_b\equiv h_bA_b$. The counterterm of $\widetilde{A}_b$ could be defined via a physical process, *e.g.* one of the decays $\widetilde{b}_1\to\widetilde{b}_2\,A$ or $A\to\widetilde{b}_1\,\widetilde{b}_2^*$, but such a definition would suffer from the problem of infrared (IR) singularities associated with gluon radiation. To overcome this problem, and given our ignorance of the

MSSM spectrum, we find less restrictive to define $\delta\widetilde{A}_b$ in terms of the $(\widetilde{b}_1\widetilde{b}_2^*A)$ proper vertex, at appropriately chosen external momenta and including suitable wave function corrections, so that the resulting combination is IR finite and gauge-independent, and gives rise to an acceptable heavy gluino limit. Denoting the proper vertex $\tilde{b}_1\tilde{b}_2^*A$ with $i\Lambda_{12A}(p_1^2,p_2^2,p_A^2)$, we define

$$\begin{aligned}\delta\widetilde{A}_b \;=\; & -\frac{i}{\sqrt{2}}\left[\Lambda_{12A}(m^2_{\tilde{b}_1},m^2_{\tilde{b}_1},0)+\Lambda_{12A}(m^2_{\tilde{b}_2},m^2_{\tilde{b}_2},0)\right] \\ & +\frac{\widetilde{A}_b}{2}\left[\frac{\Pi^{\tilde{b}}_{11}(m^2_{\tilde{b}_1})-\Pi^{\tilde{b}}_{11}(m^2_{\tilde{b}_2})}{m^2_{\tilde{b}_1}-m^2_{\tilde{b}_2}}+\frac{\Pi^{\tilde{b}}_{22}(m^2_{\tilde{b}_1})-\Pi^{\tilde{b}}_{22}(m^2_{\tilde{b}_2})}{m^2_{\tilde{b}_1}-m^2_{\tilde{b}_2}}\right. \\ & \left.+\frac{\Pi_{AA}(m^2_{\tilde{b}_1})-\Pi_{AA}(m^2_{\tilde{b}_2})}{m^2_{\tilde{b}_1}-m^2_{\tilde{b}_2}}\right]. \qquad (4)\end{aligned}$$

Although we have used Eqs. (2)–(3) to define an OS bottom Yukawa coupling h_b through the sbottom mixing, we still need to exploit the experimental information on the bottom mass in order to obtain the $\overline{\rm DR}$ running coupling $\hat{h}_b$. The OS coupling will then be computed through the relation $h_b=\hat{h}_b-\delta h_b$. We thus define the running coupling $\hat{h}_b$ at the reference scale $Q_0=175$ GeV to be

$$\hat{h}_b \equiv h_b(Q_0)^{\overline{\rm DR}}_{\rm MSSM}=\frac{\overline{m}_b\sqrt{2}}{v_1}\,\frac{1+\delta_b}{|1+\varepsilon_b|}\,, \qquad (5)$$

where: $\overline{m}_b\equiv m_b(Q_0)^{\overline{\rm DR}}_{\rm SM}=2.74\pm0.05$ GeV is the Standard Model bottom mass, evolved up to the scale Q_0 to take into account the resummation of the universal large QCD logarithms; ε_b contains the $\tan\beta$-enhanced threshold corrections from both the gluino-sbottom and the higgsino-stop loops; δ_b contains the residual threshold corrections that are not enhanced by $\tan\beta$. Notice that, as shown in Ref. [29], keeping ε_b in the denominator of Eq. (5) allows to resum the $\tan\beta$-enhanced threshold corrections to all orders in the perturbative expansion. On the other hand, there is no preferred way of including the threshold corrections parametrised by δ_b, whose effect on the value of $\hat{h}_b$ is anyway very small. It appears from Eq. (5) that $\hat{h}_b$ blows up when ε_b approaches -1, in which case the correct value of the bottom mass cannot be reproduced with $\hat{h}_b$ in the perturbative regime, and the corresponding set of MSSM parameters must be discarded.

For the top/stop sector, we take as input the physical top mass and the parameters $(m_{Q,\tilde{t}},m_U,A_t)$ that can be derived by rotating the diagonal matrix of the OS stop masses by the angle $\theta_{\tilde{t}}$, defined in Eq. (1). Concerning the sbottom sector, additional care is required, because of our non-trivial definition of h_b and of the fact that, at one loop, the parameter $m_{Q,\tilde{b}}$ entering the sbottom mass matrix differs from the corresponding stop parameter $m_{Q,\tilde{t}}$ by a finite shift [28]. We start by computing the renormalised coupling h_b as given by Eq. (3) and (5). Then we compute $m_{Q,\tilde{b}}$ following the prescription of Ref. [28]. Finally, we use the parameters h_b and $m_{Q,\tilde{b}}$ to compute the actual values of the OS sbottom masses and mixing angle.

We are now ready to discuss the numerical effect of our two-loop corrections (for details on their calculation, see Refs. [14, 15]). In Fig. 1 we show the lighter Higgs boson mass M_h as a function of $\tan\beta$ [we recall that, although our OS prescription for the sbottom sector is defined in the limit $\tan\beta \to \infty$, the corrections have an indirect dependence on $\tan\beta$ coming from the input value for $\hat{h}_b$, see Eq. (5)]. The other input parameters are

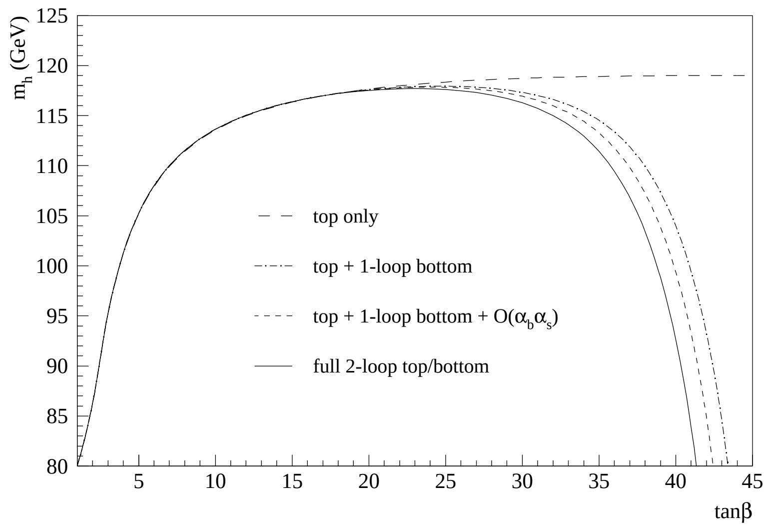

Fig. 1. Lighter Higgs boson mass M_h as a function of $\tan\beta$. The input parameters are $M_A = 120$ GeV, $A_t = 1$ TeV, $A_b = 2$ TeV, $m_{Q,\tilde{t}} = m_U = m_D = m_{\tilde{g}} = -\mu = 1$ TeV. The meaning of the different curves is explained in the text.

chosen as $M_A = 120$ GeV, $A_t = 1$ TeV, $A_b = 2$ TeV, $m_{Q,\tilde{t}} = m_U = m_D = m_{\tilde{g}} = -\mu = 1$ TeV. The long-dashed curve corresponds to the value of M_h obtained at $\mathcal{O}(\alpha_t + \alpha_t\alpha_s + \alpha_t^2)$, *i.e.* by including only the one- and two-loop corrections controlled by the top Yukawa coupling; the dot–dashed curve includes in addition the one-loop $\mathcal{O}(\alpha_b)$ corrections, controlled by the bottom Yukawa coupling; the short-dashed curve includes the two-loop $\mathcal{O}(\alpha_b\alpha_s)$ corrections computed in Ref. [14]; finally, the solid curve corresponds to the full two-loop Yukawa computation of M_h, *i.e.* it includes also the $\mathcal{O}(\alpha_t\alpha_b + \alpha_b^2)$ corrections computed in Ref. [15]. We can see from Fig. 1 that the corrections controlled by the top Yukawa coupling depend very weakly on $\tan\beta$ when the latter is large. On the other hand, the $\mathcal{O}(\alpha_b)$ corrections lower considerably M_h when $\tan\beta$ increases. Concerning the two-loop corrections controlled by the bottom Yukawa coupling, the comparison between the dot-dashed and short-dashed curves shows that the $\mathcal{O}(\alpha_b\alpha_s)$ corrections amount to a small fraction of the $\mathcal{O}(\alpha_b)$ ones, but they can still lower M_h by several GeV when $\tan\beta$ is large. The comparison between the short-dashed and solid curves shows that the effect of the $\mathcal{O}(\alpha_t\alpha_b + \alpha_b^2)$ corrections can also amount to several GeV when $\tan\beta$ is large.

It appears from Fig. 1 that the two-loop $\mathcal{O}(\alpha_b\alpha_s)$ and $\mathcal{O}(\alpha_t\alpha_b + \alpha_b^2)$ corrections are a small fraction of the one-loop $\mathcal{O}(\alpha_b)$ ones. We stress that this is a desirable consequence of our renormalisation prescription, which allows to set apart the $\tan\beta$-enhanced threshold corrections, resummed to all orders in the renormalised coupling h_b. If we were to adopt for the bottom/sbottom sector the same renormalisation prescription that we use for the top/stop sector, the dependence on $\tan\beta$ of the one-loop corrections would be smoother, but very large corrections would appear at two loops, questioning the validity of the perturbative expansion.

3. Radiative corrections in the $\overline{\mathrm{DR}}$ renormalisation scheme

In alternative to defining an OS renormalisation scheme, it is always possible to express the input parameters in a minimal subtraction scheme such as $\overline{\mathrm{DR}}$. The $\overline{\mathrm{DR}}$ renormalisation scheme is particularly convenient in constrained scenarios where the MSSM parameters at the weak scale are obtained from a set of unifying high-energy boundary conditions via suitable RGE. In the following we discuss three codes for the computation of the MSSM particle spectrum, `SoftSusy`, `SPheno` and `SuSpect`, which employ the $\overline{\mathrm{DR}}$ version of the results of Refs. [12–15] in the calculation of the MSSM Higgs boson masses. We also compare their results with those of `FeynHiggs`, which employs the OS scheme, and provide an estimate of the residual theoretical uncertainties.

3.1. Higgs boson masses in the constrained MSSM

We first work in the framework of constrained MSSM scenarios, and restrict ourselves to the case of mSUGRA, where the relevant input parameters are three universal soft SUSY-breaking terms at the high-energy scale, $m_0\,, m_{1/2}$ and A_0, the ratio of the Higgs vacuum expectation values $\tan\beta$ (expressed in the $\overline{\mathrm{DR}}$ scheme at the renormalisation scale $Q = M_Z$) and the scale-invariant sign of μ. For definiteness, we adopt the choice of parameters known as SPS1a point:

$$m_0 = 100 \text{ GeV}, \quad m_{1/2} = 250 \text{ GeV}, \quad A_0 = -100\,, \quad \tan\beta = 10\,, \quad \mu > 0\,. \tag{6}$$

We start our discussion by presenting in the first three lines of Table I the results of `SoftSusy`, `SPheno` and `SuSpect` for the physical masses of the CP-even Higgs bosons, M_h and M_H, for the physical mass of the CP-odd Higgs boson, M_A, and for the parameter μ (the latter interpreted as a $\overline{\rm DR}$ running parameter obtained by enforcing the EWSB conditions at the default renormalisation scale $M_{\rm EWSB} = \sqrt{m_{\tilde{t}_1}\, m_{\tilde{t}_2}}$).

TABLE I

Neutral Higgs boson masses and μ (in GeV) as computed by `SoftSusy`, `SPheno` and `SuSpect`. The CP-even Higgs boson masses resulting from `FeynHiggs` are also shown.

Code	M_h	M_H	M_A	μ
`SoftSusy`	112.1	406.5	406.2	364.8
`SPheno`	112.2	406.0	405.7	364.3
`SuSpect`	112.1	406.5	406.1	364.7
`FeynHiggs`	113.8	406.5	—	—

It can be seen from Table I that the agreement between the three codes is very good, the discrepancies being contained in a half GeV. For other choices of the SUSY-breaking parameters (see Ref. [16]) the discrepancies can be somewhat larger, but generally below the 1% level. We find this agreement very satisfactory, taking into account the fact that the results are obtained with three independent codes that — although based on the same set of formulae for the corrections to the Higgs boson masses and the EWSB conditions — differ in many details of the calculation. We have checked that, if we force the three codes to use the same computation of the threshold corrections to the gauge and Yukawa couplings and the same set of RGE, the residual discrepancies in the results for the neutral Higgs boson masses and μ become negligible. However, we stress that the differences among the three codes are a matter of choice, because they all correspond to effects that are of higher order with respect to the accuracy required by the calculation.

A first estimate of the residual theoretical uncertainties can be obtained by comparing the results of our $\overline{\rm DR}$ calculation of the Higgs boson masses with those of the computer code `FeynHiggs`, which employs the OS renormalisation scheme. The resulting discrepancies must be due to terms that are formally of higher order, *i.e.* two-loop terms controlled by the electroweak gauge couplings and three-loop terms, among which the most important are controlled by the top Yukawa and strong gauge couplings. As `FeynHiggs` does not perform the evolution of the MSSM parameters from the high-energy input scale to $M_{\rm EWSB}$, the input parameters at the weak scale (including M_A and μ) are taken from the output of `SuSpect`. The last line of Table I shows the results of `FeynHiggs` for M_h and M_H: it can be seen that the difference with `SuSpect` in the value of M_h is of the order of 2 GeV (the same occurs for other typical scenarios of constrained MSSM). The excellent agreement in the value of M_H is due to the strict correlation between the latter and the value of M_A, which is taken from the output of `SuSpect`.

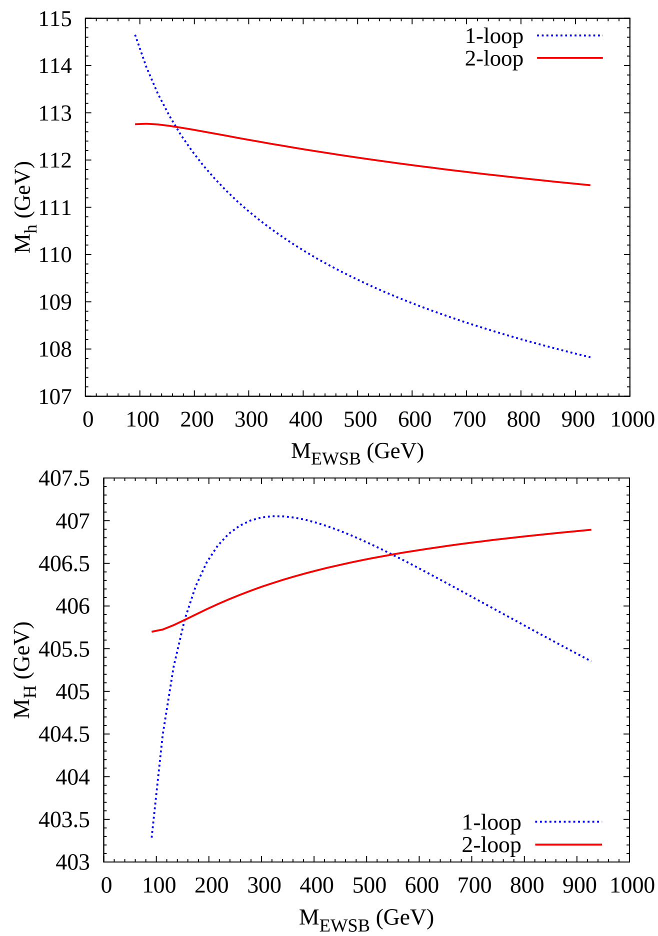

Fig. 2. The CP-even Higgs boson masses, M_h (top) and M_H (bottom), as a function of the minimisation scale of the effective potential, $M_{\rm EWSB}$.

Another measure of the effect of the higher orders consists in studying the numerical dependence of the results for the physical Higgs boson masses on the renormalisation scale $M_{\rm EWSB}$ at which the effective potential is minimised and the radiatively corrected masses are computed. In the ideal case of an all-orders calculation, the physical observables should not depend on the choice of the scale. The residual scale dependence still present in the real case can be taken as a rough estimate of the magnitude of the corrections that are left uncomputed.

The top and bottom panels of Fig. 2 show the renormalisation scale dependence of the CP-even Higgs boson masses M_h and M_H, respectively, as computed by `SoftSusy` in the SPS1a scenario. The dotted line in each plot corresponds to the one-loop computation of the relevant mass, whereas the solid line corresponds to the two-loop computation. It can be seen from the top panel of Fig. 2 that the one-loop results for the lighter Higgs boson mass M_h show a sizeable scale dependence, varying by nearly 8 GeV in the considered range of $M_{\rm EWSB}$. On the other hand, it appears that the inclusion of the two-loop corrections significantly improves the scale dependence of M_h, leaving a residual variation of the order of 2 GeV in the considered range of $M_{\rm EWSB}$. The bottom panel of Fig. 2 shows that a similar situation occurs in the case of the heavier Higgs boson mass M_H, with the inclusion of the two-loop corrections clearly improving the renormalisation scale dependence. Again, similar results can be obtained for other typical scenarios of constrained MSSM.

3.2. Higgs boson masses in the unconstrained MSSM

In alternative to setting the input parameters at the high-energy scale, `SPheno` and `SuSpect` allow the user to set arbitrarily the MSSM input parameters at the weak scale as $\overline{\rm DR}$-renormalised quantities (the same option will soon be implemented in `SoftSusy`). In Table II we show the values of the lighter Higgs boson mass as obtained by `SPheno`, `SuSpect` and, for comparison, `FeynHiggs`, for three different values of the stop mixing parameter $X_t = A_t - \mu \cot\beta$. The other relevant MSSM parameters, *i.e.* a common stop mass term M_S, the gluino mass M_3, μ and M_A are set to 1 TeV, and $\tan\beta = 10$.

It can be seen from Table II that the results of `SPheno` and `SuSpect` agree well, the discrepancies being of the order of half GeV. The third line of the table shows that for small or moderate stop mixing the discrepancy with the OS calculation of `FeynHiggs` is of the order of 2 GeV, as was the case in the constrained MSSM (see Table I). However, for large stop mixing the difference between the $\overline{\rm DR}$ and OS computations reaches 4–5 GeV, indicating that in this case the uncertainty due to higher order corrections is larger.

TABLE II

Lighter Higgs boson mass M_h (in GeV) in the unconstrained MSSM as computed by `SPheno`, `SuSpect` and `FeynHiggs` for three values of the stop mixing parameter X_t.

Code	$X_t = 0$	$X_t = M_S$	$X_t = \sqrt{6}M_S$
`SPheno`	114.3	118.8	130.0
`SuSpect`	113.8	118.4	129.4
`FeynHiggs`	115.4	120.3	133.8

This work was partially supported by the European Community's Human Potential Programme HPRN-CT-2000-00149 (Collider Physics).

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