SUSY HIGGS BOSON DECAYS

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ABSTRACT

I discuss the decay modes of the neutral and charged Higgs bosons in the minimal supersymmetric extension of the Standard Model. Special emphasis will be put on the QCD corrections to the hadronic decay modes, the below threshold –three body– decays and the decays into supersymmetric particles, charginos, neutralinos and sfermions. A Fortran code calculating the various Higgs decay branching ratios is then briefly presented.

1. Introduction

In the Minimal Supersymmetric extension of the Standard Model (MSSM), the Higgs sector is extended to comprise three neutral \( h/H \) (CP=+), \( A \) (CP=–) and a pair of charged scalar particles \( H^\pm \). The Higgs sector is highly constrained since there are only two free parameters at tree–level: a Higgs mass parameter \( M_A \) and the ratio of the vacuum expectation values of the two doublet fields responsible for the symmetry breaking, \( \tan \beta \) [which in Grand Unified Supersymmetric models with \( b–\tau \) Yukawa coupling unification is forced to be either small, \( \tan \beta \sim 1.5 \), or large, \( \tan \beta \sim 50 \)]. After the inclusion of the large radiative corrections, while the lightest Higgs boson \( h \) is predicted to be lighter than \( M_h \lesssim 130 \) GeV, the \( H, A \) and \( H^\pm \) states are expected to have masses of the order of a few hundred GeV.

The decay pattern of the MSSM Higgs bosons is determined to a large extent by their couplings to fermions and gauge bosons, which in general depend strongly on \( \tan \beta \) and the mixing angle \( \alpha \) in the CP–even sector. The pseudoscalar and charged Higgs boson couplings to down (up) type fermions are (inversely) proportional to \( \tan \beta \); the pseudoscalar \( A \) has no tree level couplings to gauge bosons. For the CP–even Higgs bosons, the couplings to down (up) type fermions are enhanced (suppressed) compared to the SM Higgs couplings \( [\tan \beta > 1] \); the couplings to gauge bosons are suppressed by \( \sin \alpha / \cos (\beta – \alpha) \) factors [see Table 1.]

For large values of \( \tan \beta \) the pattern is simple, a result of the strong enhancement of the Higgs couplings to down–type fermions. The neutral Higgs bosons will decay into \( b\bar{b} \) (\( \sim 90\% \)) and \( \tau^+\tau^- \) (\( \sim 10\% \)) pairs, and \( H^\pm \) into \( \tau \nu_\tau \) pairs below and \( tb \) pairs above the top–bottom threshold. For the CP–even Higgs bosons \( h \) and \( H \), only when \( M_h \)

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approaches its maximal value is this simple rule modified: in this decoupling limit, the $h$ boson is SM–like and decays into charm and gluons with a rate similar to the one for $\tau^+\tau^− \sim 5\%$ and in the high mass range, $M_h \sim 130$ GeV, into $W$ pairs with one of the $W$ bosons being virtual; the $H$ boson will mainly decay into $hh$ and $AA$ final states.

$\Phi \rightarrow Q\overline{Q}$ = $3G_FM_\Phi^2g_{\Phi QQ}^2m_\Phi^2(M_\Phi) \left[ 1 + 5.67\frac{\alpha_s}{\pi} + (35.94 - 1.36N_F)\frac{\alpha_s^2}{\pi^2} \right]$(1)

in the $\overline{MS}$ renormalization scheme; the running quark mass and the QCD coupling are defined at the scale of the Higgs mass, absorbing this way any large logarithms.

2. Hadronic Decay Modes: QCD Corrections

The particle width for decays to massless $b, c$ quarks directly coupled to the Higgs particle is given, up to $O(\alpha_s^2)$ QCD corrections [the effect of the electroweak radiative corrections in the branching ratios is negligible], by the well-known expression

$\Gamma[\Phi \rightarrow Q\overline{Q}] = \frac{3G_FM_\Phi^2}{4\sqrt{2}\pi}g_{\Phi QQ}^2m_\Phi^2(M_\Phi) \left[ 1 + 5.67\frac{\alpha_s}{\pi} + (35.94 - 1.36N_F)\frac{\alpha_s^2}{\pi^2} \right]$(1)

For small values of $\tan\beta \sim 1$ the decay pattern of the heavy neutral Higgs bosons is much more complicated. The $b$ decays are in general not dominant any more; instead, cascade decays to pairs of light Higgs bosons and mixed pairs of Higgs and gauge bosons are important and decays to $WW/ZZ$ pairs will play a role. For very large masses, they decay almost exclusively to top quark pairs. The decay pattern of the charged Higgs bosons for small $\tan\beta$ is similar to that at large $\tan\beta$ except in the intermediate mass range where cascade decays to $Wh$ are dominant.

When the decays into supersymmetric particles are kinematically allowed, as it should be the case at least for the heavy CP–even, CP–odd and charged Higgs bosons, the pattern becomes even more complicated since the decay channels into charginos, neutralinos and squarks play will play a non–negligible role.

In the following, I will discuss three topics related to the decay modes of the Higgs particles in the MSSM: (a) the QCD corrections to the hadronic decay modes, (b) the below threshold three–body decays and (c) the decays into SUSY particles of the heavy $H, A$ and $H^\pm$ bosons, including the QCD corrections to the squark decay modes. I will then briefly introduce a Fortran code which calculates the various decay branching ratios. For more details and for a complete list of references, see the original papers Refs. [5–9].

Table 1: Higgs couplings to fermions and gauge bosons normalized to the SM Higgs couplings, and their limit for $M_A \gg M_Z$.

| $\Phi$ | $g_{\Phi uu}$ | $g_{\Phi dd}$ | $g_{\Phi vv}$ |
|-------|-------------|-------------|-------------|
| $h$   | $\cos \alpha / \sin \beta \rightarrow 1$ | $-\sin \alpha / \cos \beta \rightarrow 1$ | $\sin(\beta - \alpha) \rightarrow 1$ |
| $H$   | $\sin \alpha / \sin \beta \rightarrow 1/tg\beta$ | $\cos \alpha / \cos \beta \rightarrow tg\beta$ | $\cos(\beta - \alpha) \rightarrow 0$ |
| $A$   | $1/tg\beta$ | $tg\beta$ | 0 |
The quark masses can be neglected in general except for top quark decays where this approximation holds only sufficiently far above threshold. Since the relation between the charm pole mass $M_c$ and the $\overline{\text{MS}}$ mass evaluated at the pole mass $\overline{m}_c(M_c)$ is badly convergent, one can adopt the running quark masses $\overline{m}_Q(M_Q)$ [which have been extracted directly from QCD sum rules evaluated in a consistent $\mathcal{O}(\alpha_s)$ expansion] as starting points. The evolution from $M_Q$ to a scale $\mu \sim M_{\Phi}$ is given by:

$$\overline{m}_Q(\mu) = \overline{m}_Q(M_Q) c[\alpha_s(\mu)/\pi]/c[\alpha_s(M_Q)/\pi]$$

$$c(x) = (25/6x)^{12/25} [1 + 1.014x + 1.39x^2] \quad \text{for} \quad M_c < \mu < M_b$$

$$c(x) = (23/6x)^{12/23} [1 + 1.175x + 1.50x^2] \quad \text{for} \quad M_b < \mu$$

(2)

Typical values of the running $b, c$ masses at the scale $\mu = 100$ GeV, characteristic for $M_{\Phi}$, are displayed in Table 2, with the evolution calculated for $\alpha_s(M_Z) = 0.118 \pm 0.006$; $M_Q^{\alpha_2}$ are the quark pole masses.

| $\alpha_s(M_Z)$ | $\overline{m}_Q(M_Q)$ | $M_Q = M_Q^{\alpha_2}$ | $\overline{m}_Q(\mu = 100 \text{ GeV})$ |
|-----------------|----------------------|------------------------|---------------------------------|
| $b$ 0.112        | (4.26 ± 0.02) GeV    | (4.62 ± 0.02) GeV      | (3.04 ± 0.02) GeV               |
| 0.118           | (4.23 ± 0.02) GeV    | (4.62 ± 0.02) GeV      | (2.92 ± 0.02) GeV               |
| 0.124           | (4.19 ± 0.02) GeV    | (4.62 ± 0.02) GeV      | (2.80 ± 0.02) GeV               |
| $c$ 0.112        | (1.25 ± 0.03) GeV    | (1.42 ± 0.03) GeV      | (0.69 ± 0.02) GeV               |
| 0.118           | (1.23 ± 0.03) GeV    | (1.42 ± 0.03) GeV      | (0.62 ± 0.02) GeV               |
| 0.124           | (1.19 ± 0.03) GeV    | (1.42 ± 0.03) GeV      | (0.53 ± 0.02) GeV               |

Table 2: The running $b$ and $c$ quark masses in the $\overline{\text{MS}}$ scheme at a scale $\mu = 100$ GeV.

The decay of the Higgs bosons to gluons is, to a good approximation mediated by heavy top quark loops; the partial decay width, including QCD radiative corrections which are built up by the exchange of virtual gluons and the splitting of a gluon into two gluons or into $N_F$ massless quark–antiquark pairs, is given by $[\mu \sim M_{\Phi}]$

$$\Gamma^{N_F}[\phi \rightarrow gg + ..] = \frac{G_F g_{\phi t}^2 a_s^2 M_{\phi}^3}{36\sqrt{2}\pi^3} \left[ 1 + \frac{\alpha_s}{\pi} \left( \frac{95}{4} - \frac{7}{6}N_F + \frac{33 - 2N_F}{6} \log \frac{\mu^2}{M_{\phi}^2} \right) \right]$$

$$\Gamma^{N_F}[A \rightarrow gg + ..] = \frac{G_F g_{A t}^2 a_s^2 M_A^3}{16\sqrt{2}\pi^3} \left[ 1 + \frac{\alpha_s}{\pi} \left( \frac{97}{4} - \frac{7}{6}N_F + \frac{33 - 2N_F}{6} \log \frac{\mu^2}{M_A^2} \right) \right]$$

(3)

with $\phi = h, H$ and $\alpha_s \equiv \alpha_s^{N_F}(\mu^2)$. The radiative corrections are very large, nearly doubling the partial width. The final states $\Phi \rightarrow b\bar{b}g$ and $c\bar{c}g$ are also generated through processes in which the $b, c$ quarks are coupled to the Higgs boson directly. Gluon splitting $g \rightarrow b\bar{b}$ in $\Phi \rightarrow gg$ increases the inclusive decay probabilities $\Gamma(\Phi \rightarrow b\bar{b} + ..)$ etc. Since $b$ quarks, and eventually $c$ quarks, can in principle be tagged experimentally, it is physically meaningful to consider the particle width of Higgs.
decays to gluon and light $u,d,s$ quark final jets separately. The contribution of $b,c$ quark final states to the coefficient in front of $\alpha_s$ in eq. (3) is:

$$-\frac{7}{3} + \frac{1}{3} \left[ \log \frac{M_b^2}{M_\Phi^2} + \log \frac{M_c^2}{M_\Phi^2} \right]$$

Instead of naively subtracting this contribution, it may be noticed that the mass logarithms can be absorbed by changing the number of active flavors from $N_F = 5$ to $N_F = 3$ in the QCD coupling $\alpha_s^{(N_F)}$. The subtracted parts may be added to the partial decay widths into $c$ and $b$ quarks.

The numerical analysis of the branching ratios for the lightest CP–even Higgs decays in the decoupling limit where $h$ is SM–like, with the quark masses and QCD couplings given above and a top mass $M_t = (176 \pm 11)$ GeV, is shown in Fig. 1. To estimate systematic uncertainties, the variation of the $c$ mass has been stretched over $2\sigma$ and the uncertainty of the $b$ mass to 0.05 GeV. However, the dominant error in the predictions is due to the uncertainty in $\alpha_s$ and the errors in the prediction for the charm and gluon branching ratios are very large. Nevertheless, the expected hierarchy of the Higgs decay modes is clearly visible despite these uncertainties. Similar results hold for the heavy CP–even and CP–odd Higgs decays.

![Fig. 1: Branching ratios of the $h$ boson in the decoupling limit, including the uncertainties from the quark masses and the QCD coupling $\alpha_s$ [shaded bands].](image)
3. Three Body decay modes

Besides these two–body decays, below–threshold modes can play an important role. It is well–known that SM Higgs decays into real and virtual Z pairs are quite substantial: the suppression by the off–shell propagator and the additional $Z f f$ coupling is at least partly compensated by the large Higgs coupling to the Z bosons. For the same reason, three–body decays of MSSM Higgs particles mediated by gauge bosons, heavy Higgs bosons and top quarks, are of physical interest. Important three–body decays for the $H, A$ and $H^\pm$ bosons are $[V = W, Z]$:

$$H \rightarrow V V^* \rightarrow V f f^{(*)}, \quad AZ^* \rightarrow A f f^*, \quad H^\pm W^{\mp*} \rightarrow H^\pm f f^*, \quad t\bar{t}^* \rightarrow t\bar{b}W^+$$ (4)

$$A \rightarrow hZ^* \rightarrow h f f^*, \quad t\bar{t}^* \rightarrow t\bar{b}W^+$$ (5)

$$H^\pm \rightarrow hW^* \rightarrow h f f^*, \quad AW^* \rightarrow A f f^*, \quad b\bar{t}^* \rightarrow b\bar{b}W$$ (6)

For the lightest Higgs boson $h$, the only releveant below threshold decay mode is $h \rightarrow W^*W^*$ for $M_h \sim 130$ GeV. In this case, both the $W$’s have to be taken off–shell. The branching ratios for $h, H, A$ and $H^\pm$ decays are shown in Fig. 2 for $\tan \beta = 1.5$, in the case where the mixing in the stop sector is neglected.

For the heavy Higgs boson $H$, the decay $H \rightarrow hh$ is the dominant channel, superseded by $t\bar{t}$ decays above the threshold [for the latter, the inclusion of the three–body modes provides a smooth transition from below to above threshold]. This rule is only broken for Higgs masses of about 140 GeV where an accidentally small value of the $\lambda_{hhh}$ coupling allows the $b\bar{b}$ and $WW^*$ decay modes to become dominant. Important channels in general, below the $t\bar{t}$ threshold, are decays to pairs of gauge bosons and $b\bar{b}$ decays. In a restricted range of $M_H$, below–threshold $AZ^*$ and $H^\pm W^{\mp*}$ also play a non–negligible role. In the case of the pseudoscalar $A$, the dominant modes are the $A \rightarrow b\bar{b}$ and $A \rightarrow t\bar{t}$ decays below the $hZ$ and $t\bar{t}$ thresholds respectively; in the intermediate mass region, $M_A = 200$ to 300 GeV, the decay $A \rightarrow hZ^*$ [which reaches $\sim 1\%$ already at $M_A = 130$ GeV] dominates. The gluonic decays are significant around the $t\bar{t}$ threshold. For the charged Higgs boson, the inclusion of the three–body decay modes will reduce the branching ratio for the $\tau\nu$ channel quite significantly. Indeed, this decay does not overwhelm all the other modes since the three–body decay channels $H^+ \rightarrow hW^*$ as well as $H^+ \rightarrow AW^*$ in the low mass range and $H^+ \rightarrow b\bar{t}^*$ in the intermediate mass range have appreciable branching ratios.

The total widths of the Higgs bosons are in general considerably smaller than for the SM Higgs due to the absence or the suppression of the decays to $W/Z$ bosons which grow as $M_H^3$. The dominant decays are built-up by top quarks so that the widths rise only linearly with $M_\Phi$. However, for large $\tan \beta$ values, the decay widths scale in general like $\tan^2 \beta$ and can become experimentally significant, for $\tan \beta \gtrsim \mathcal{O}(30)$ and for large $M_\Phi$. 

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Fig. 2: Branching ratios for the CP–even, the CP–odd and the charged MSSM Higgs bosons, including the three–body decays, for $\tan \beta = 1.5$ and no stop mixing.
4. SUSY Decay modes

In the previous discussion, we have assumed that decay channels into neutralinos, charginos and sfermions are shut. However, these channels could play a significant role, since some of these particles can have masses in the $O(100 \text{ GeV})$ range or less. To discuss these decays, we will restrict ourselves to the MSSM constrained by minimal Supergravity, in which the SUSY sector is described in terms of five universal parameters at the GUT scale: the common scalar mass $m_0$, the common gaugino mass $M_{1/2}$, the trilinear coupling $A$, the bilinear coupling $B$ and the higgsino mass $\mu$. These parameters evolve according to the RGEs, forming the supersymmetric particle spectrum at low energy.

The requirement of radiative electroweak symmetry breaking further constrains the SUSY spectrum, since the minimization of the one–loop Higgs potential specifies the parameter $\mu$ [to within a sign] and also $B$. The unification of the $b$ and $\tau$ Yukawa couplings gives another constraint: in the $\lambda_t$ fixed–point region, the value of $\tan\beta$ is fixed by the top quark mass through: $m_t \simeq (200 \text{ GeV}) \sin\beta$, leading to $\tan\beta \simeq 1.75$. There also exists a high–$\tan\beta$ [$\lambda_b$ and $\lambda_\tau$ fixed–point] region for which $\tan\beta \sim 50$–60. If one also notes that moderate values of the trilinear coupling $A$ have little effect on the resulting spectrum, then the whole SUSY spectrum will be a function of $\tan\beta$ which we take to be $\tan\beta = 1.75$ and 50, the sign of $\mu$, $m_0$ which in practice we replace with $M_A$ taking the two illustrative values $M_A = 300$ and 600 GeV, and the common gaugino mass $M_{1/2}$ that we will freely vary.

The decay widths of the heavy CP-even, the CP–odd and the charged Higgs bosons, $H, A$ and $H^\pm$, into pairs of neutralinos and charginos [dashed lines], squarks [long–dashed lines] and sleptons [dot–dashed lines], as well as the total [solid lines] and non–SUSY [dotted–lines] decay widths, are shown in Fig. 3 for $\tan\beta = 1.75$, $\mu > 0$ and two values of $M_A = 300$ [left curves] and 600 GeV [right curves].

For $M_A = 300$ GeV, i.e. below the $t\bar{t}$ threshold, the widths of the $H$ decays into inos and sfermions are much larger than the non–SUSY decays. In particular, squark [in fact $\tilde{t}$ and $\tilde{b}$ only] decays are almost two–orders of magnitude larger when kinematically allowed. The situation changes dramatically for larger $M_A$ when the $t\bar{t}$ channel opens up: only the decays into $\tilde{t}$ pairs when allowed are competitive with the dominant $H \rightarrow t\bar{t}$ channel. Nevertheless, the decays into inos are still substantial having BRs at the level of 20%; the decays into sleptons never exceed a few percent.

In the case of the pseudoscalar $A$, because of CP–invariance and the fact that sfermion mixing is small except in the stop sector, only the decays into inos and $A \rightarrow \tilde{t}_1 \tilde{t}_2$ decays are allowed. For these channels, the situation is quite similar to the case of $H$: below the $t\bar{t}$ threshold the decay width into ino pairs is much larger than the non–SUSY decay widths [here $\tilde{t}_2$ is too heavy for the $A \rightarrow \tilde{t}_1 \tilde{t}_2$ decay to be
allowed], but above $2m_t$ only the $A \rightarrow \tilde{t}_1 \tilde{t}_2$ channel competes with the $t\bar{t}$ decays.

For the charged Higgs boson $H^\pm$, only the decay $H^+ \rightarrow \tilde{t}_1 \tilde{b}_1$ [when kinematically allowed] competes with the dominant $H^+ \rightarrow t\bar{b}$ mode, yet the $\tilde{\chi}^+\tilde{\chi}^0$ decays have a branching ratio of a few ten percent; the decays into sleptons are at most of $\mathcal{O}(\%)$.

In the case where $\mu < 0$, the situation is quite similar as above. For large $\text{tg} \beta$ values, $\text{tg} \beta \sim 50$, all gauginos and sfermions are very heavy and therefore kinematically inaccessible, except for the lightest neutralino and the $\tau$ slepton. Moreover, the $b\bar{b}/\tau\tau$ and $t\bar{b}/\tau\nu$ [for the neutral and charged Higgs bosons respectively] are enhanced so strongly, that they leave no chance for the SUSY decay modes to be significant. Therefore, for large $\text{tg} \beta$, the simple pattern of $b\bar{b}/\tau\tau$ and $tb$ decays for heavy neutral and charged Higgs bosons still holds true even when the SUSY decays are allowed.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig3.png}
\caption{Decay widths for the SUSY decay modes of the heavy CP–even, CP–odd and charged Higgs bosons, for $\text{tg} \beta = 1.75$. The total and the non–SUSY widths are also shown.}
\end{figure}
Since the decays into stop and sbottom squarks can be dominant when kinematically allowed, QCD corrections must be incorporated in order to have full control on the decay widths and to make a reliable comparison with the standard (non–SUSY) decay channels. The QCD corrections to the decays of the heavy CP–even, CP–odd and charged MSSM Higgs bosons into stop and sbottom quarks have been recently calculated\(^8\). These corrections are found to be rather large, enhancing or suppressing the widths by amounts up to 50% and in some case even more. The QCD corrections depend strongly on the gluino mass; however for very heavy gluinos, they are only logarithmically dependent on \(m_{\tilde{g}}\). Contrary to the case of Higgs boson decays into light quark pairs, these large corrections cannot be absorbed into running squark masses since the latter are expected to be of the same order as as the Higgs masses.

5. The program HDECAY\(^9\)

Finally, let me shortly describe the fortran code HDECAY, which calculates the various decay widths and the branching ratios of Higgs bosons in the SM and the MSSM and which includes:

(a) All decay channels that are kinematically allowed and which have branching ratios larger than \(10^{-4}\), \(y\) compris the loop mediated, the three body decay modes and in the MSSM the cascade and the supersymmetric decay channels.
(b) All relevant two-loop QCD corrections to the decays into quark pairs and to the quark loop mediated decays into gluons are incorporated in the most complete form; the small leading electroweak corrections are also included.
(c) Double off–shell decays of the CP–even Higgs bosons into massive gauge bosons which then decay into four massless fermions, and all all important below–threshold three–body decays discussed previously.
(d) In the MSSM, the complete radiative corrections in the effective potential approach with full mixing in the stop/sbottom sectors; it uses the renormalisation group improved values of the Higgs masses and couplings and the relevant leading next–to–leading–order corrections are also implemented.
(e) In the MSSM, all the decays into SUSY particles (neutralinos, charginos, sleptons and squarks including mixing in the stop, sbottom and stau sectors) when they are kinematically allowed. The SUSY particles are also included in the loop mediated \(\gamma\gamma\) and \(gg\) decay channels.

The basic input parameters, fermion and gauge boson masses and total widths, coupling constants and in the MSSM, soft–SUSY breaking parameters can be chosen from an input file. In this file several flags allow to switch on/off or change some options [\(e.g.\) chose a particular Higgs boson, include/exclude the multi–body or SUSY decays, or include/exclude specific higher–order QCD corrections]. The results for the many decay branching ratios and the total decay widths are written to several output files with headers indicating the processes and giving the input parameters.
6. Acknowledgments:

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7. References

1. For a review see, H. Haber and G. Kane, Phys. Rep. 117 (1985) 75.
2. For a review, see J.F. Gunion, H.E. Haber, G.L. Kane and S. Dawson, *The Higgs Hunter’s Guide*, Addison–Wesley, Reading 1990.
3. For a recent summary, see M. Carena and P. Zerwas [conv.] et al. *Higgs Physics at LEPII*, CERN yellow report CERN-96-01, edited by G. Altarelli, T. Sjostrand and F. Zwirner.
4. Y. Okada, M. Yamaguchi and T. Yanagida, Prog. Theor. Phys. 85 (1991) 1; H. Haber and R. Hempfling, Phys. Rev. Lett. 66 (1991) 1815; J. Ellis, G. Ridolfi and F. Zwirner, Phys. Lett. 257B (1991) 83; R. Barbieri, F. Caravaglios and M. Frigeni, Phys. Lett. 258B (1991) 167.
5. A. Djouadi, M. Spira and P. Zerwas, Z. Phys. C70 (1996) 427.
6. A. Djouadi, J. Kalinowski and P. Zerwas, Z. Phys. C70 (1996) 435.
7. A. Djouadi, J. Kalinowski, P. Ohmann and P. Zerwas, hep-ph/9605339 (Z. Phys. C to appear); A. Djouadi, P. Janot, J. Kalinowski and P. Zerwas, Phys. Lett. B376 (1996) 220.
8. A. Bartl et al., hep-ph/9701398 A. Arhrib, A. Djouadi, W. Hollik and C. Jünger, hep-ph/9702426.
9. A. Djouadi, J. Kalinowski and M. Spira, PM–97–02, to appear. The program can obtained by sending an E-mail to one of the authors: djouadi@lpm.univ-montp2.fr, kalino@x4u2.desy.de or spira@cern.ch or directly from the WWW at http://www.lpm.univ-montp2.fr/~djouadi/program.html or http://wwwcn.cern.ch/~mspira/.