Three–body decays of Higgs bosons at LEP2 and application to a hidden fermiophobic Higgs.

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Abstract

We study the decays of Higgs bosons to a lighter Higgs boson and a virtual gauge boson in the context of the non–supersymmetric Two–Higgs–Doublet–Model (2HDM). We consider the phenomenological impact at LEP2 and find that such decays, when open, may be dominant in regions of parameter space and thus affect current Higgs boson search techniques. Three–body decays would be a way of producing light neutral Higgs bosons which have so far escaped detection at LEP due to suppressed couplings to the $Z$, and are of particular importance in the 2HDM (Model I) which allows both a light fermiophobic Higgs and a light charged scalar.

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1 Introduction

The Higgs sector of the Standard Model (SM) is still experimentally untested, and so far only a lower bound on the mass of the Higgs boson ($M_{\phi^0} \geq 87.6$ GeV) has been obtained. The minimal SM possesses one complex scalar doublet with a non-zero vacuum expectation value (VEV), and after symmetry breaking predicts a neutral Higgs boson. Enlarged Higgs sectors with $N$ doublets may be considered, and predict charged Higgs bosons ($H^\pm$) and additional neutral scalars. Accurate predictions of the branching ratios of these particles are needed in order to facilitate the searches at future colliders, and the present work considers decays of a Higgs boson to a lighter Higgs boson and a virtual vector boson. We shall be focusing on the Two–Higgs–Doublet–Model (2HDM), and how the presence of these three–body decays may affect the current search techniques. Some attention will be also given to a general model with $N \geq 3$ doublets, which we shall call the Multi–Higgs–Doublet–Model (MHDM). One particular form of the 2HDM has received substantial attention in the literature, mainly due to the fact that it is the structure of the minimal supersymmetric extension (MSSM) of the SM. However, there are four variants of the 2HDM which differ in how the doublets are coupled to the fermions (we are concerned with natural flavour conservation). In Ref. these are referred to as Models I, I’, II and II’, with Model II appropriate for the MSSM. The phenomenology of the four models can be quite different, both in the charged and neutral sector.

Model I has received relatively little attention in the literature, although among other features allows the possibility of a $H^\pm$ in the range of LEP2 and the phenomena known as "fermiophobia". Fermiophobic Higgs bosons ($H_F$) are searched for actively at the Tevatron and LEP, using direct production methods that make use of the $ZZH_F$ coupling. Existing limits ($M_F \geq 90$ GeV, 95% c.l) only apply for a $H_F$ with SM strength coupling, although in general this coupling will be suppressed, thus allowing a lighter $H_F$ to be hidden. This suppression is always possible in the general 2HDM for the lighter CP–even eigenstate $h$, and allows the possibility of an undetected Higgs boson with $M_h \leq 40$ GeV. Such a hidden Higgs boson, whether fermiophobic or not, could be produced by the above mentioned 3–body decay of a heavier Higgs boson if the branching ratio (BR) were sufficiently large. If the BR were dominant, present Higgs search techniques in these models would need to be changed. The aforementioned three–body decays have so far only appeared in the context of the MSSM, and have limited importance, although we shall see that their strength can be considerably greater in the general 2HDM.

The paper is organized as follows. In Section 2 we introduce the models in question and display the couplings of the Higgs bosons to the fermions. Section 3 considers constraints on the masses and couplings of the scalars from precision measurements. Section 4 briefly reviews the current literature on light, hidden neutral Higgs bosons, while Section 5 considers the impact of the three–body decay channels on the BRs.
In Section 6 we build on the work of Ref. [7] and consider the possibility of a large BR ($H^± \rightarrow cb$) which is only allowed in the MHDM. Finally, Section 7 contains our conclusions.

2 The Models

The theoretical structure of the 2HDM is well known [4], while the charged Higgs sector and neutral Higgs sectors of the MHDM have been studied in Ref. [5] and Ref. [8] respectively. The CP conserving 2HDM which is usually considered in the literature contains an important parameter

$$\tan \beta = \frac{v_2}{v_1}$$

(1)

with $v_1$ and $v_2$ being real vacuum expectation values (VEVs) of the two Higgs doublets, and $v^2 = \sum_{i=1}^{N} v_i^2 = 246^2 \text{ GeV}^2$ for $N$ doublets. In a MHDM it is usually assumed that one of the charged scalars is much lighter than the others and thus dominates the low–energy phenomenology. For the charged Higgs interactions with the fermions the relevant part of the Lagrangian is

$$\mathcal{L} = (2\sqrt{2}G_F)(XU_LV M_D D_R + YU_R V M_U D_L + Z\bar{N}_L M_E E_R)H^+ + h.c.$$  (2)

Here $U_L, U_R (D_L, D_R)$ denote left– and right–handed up (down) type quark fields, $N_L$ is the left–handed neutrino field, and $E_R$ the right–handed charged lepton field. $M_D, M_U, M_E$ are the diagonal mass matrices of the down type quarks, up type quarks and charged leptons respectively. $V$ is the CKM matrix. For the four distinct versions of the 2HDM the couplings $X, Y$ and $Z$ are given by the entries in Table 1 [6]. In the

|         | Model I | Model I' | Model II | Model II' |
|---------|---------|----------|----------|-----------|
| $X$     | $-\cot \beta$ | $-\cot \beta$ | $\tan \beta$ | $\tan \beta$ |
| $Y$     | $\cot \beta$ | $\cot \beta$ | $\cot \beta$ | $\cot \beta$ |
| $Z$     | $-\cot \beta$ | $\tan \beta$ | $\tan \beta$ | $-\cot \beta$ |

Table 1: The values of $X, Y$ and $Z$ in the 2HDM

MHDM $X, Y$ and $Z$ are arbitrary complex numbers which originate from a $N \times N$ charged scalar mass matrix. It is apparent that the models may differ significantly in their phenomenology. For the couplings of the CP–odd pseudoscalar ($A$), one may use Table 1 with $X$ interpreted as the coupling to $\bar{d}d$, $Y$ the coupling to $\bar{u}u$ and $Z$ the coupling to $\bar{f}f$. For the lighter CP–even eigenstate $h$ one finds the values given in Table 2, with $\alpha$ a mixing angle in the CP–even sector. For the heavier CP–even $H$ one must make the replacements $\cos \alpha \rightarrow \sin \alpha$ and $-\sin \alpha \rightarrow \cos \alpha$ in Table 2.
Table 2: The fermion couplings of $h$ relative to those for the minimal SM Higgs boson ($\phi^0$).

\begin{tabular}{|c|c|c|c|c|}
\hline
 & Model I & Model I' & Model II & Model II' \\
\hline
$hu\bar{u}$ & $\cos \alpha/\sin \beta$ & $\cos \alpha/\sin \beta$ & $\cos \alpha/\sin \beta$ & $\cos \alpha/\sin \beta$ \\
$hd\bar{d}$ & $\cos \alpha/\sin \beta$ & $\cos \alpha/\sin \beta$ & $- \sin \alpha/\cos \beta$ & $- \sin \alpha/\cos \beta$ \\
$hl\bar{l}$ & $\cos \alpha/\sin \beta$ & $- \sin \alpha/\cos \beta$ & $- \sin \alpha/\cos \beta$ & $\cos \alpha/\sin \beta$ \\
\hline
\end{tabular}

3 Constraints and Branching Ratios

Precision measurements of the process $b \to s\gamma$ impose the severest constraints on the mass of the charged scalar of the 2HDM (Models II and II'). For a general review of how new physics affects this decay see Ref. [11]. The diagrams which contribute to this process are essentially the same as those for the SM with the $W^\pm$ replaced by $H^\pm$.

The CLEO collaboration obtained the value [20]
\[
\text{BR}(b \to s\gamma) = (2.32 \pm 0.51 \pm 0.29 \pm 0.32) \times 10^{-4},
\]
and more recently ALEPH [21] have found
\[
\text{BR}(b \to s\gamma) = (3.38 \pm 0.74 \pm 0.85) \times 10^{-4}.
\]

It is known that for $H^\pm$ of Model I, I' and MHDM one cannot obtain a mass bound independent of $X$ and $Y$ (e.g. see Refs. [5], [7]). For $H^\pm$ of Model II and II' one finds the constraint $M_{H^\pm} \geq 330$ GeV for any value of $\tan \beta$ [22]. Measurements of $b \to s\gamma$ do constrain $\tan \beta$ in Model I and I' for a given Higgs mass although the decay $Z \to b\bar{b}$ imposes stronger constraints. Ref. [22] shows that from the latest $R_b$ measurements one can obtain the respective bounds of $\tan \beta \geq 1.8, 1.4, 1.0$ (95% c.l) for $M_{H^\pm} = 85, 200, 425$ GeV. The BRs for $H^\pm$ of mass 80 GeV are given in Table 3, excluding the possibility of three–body decays. The Higgs mass determines the energy scale of the decay and so one must evaluate the quark masses at the scale $Q = M_{H^\pm}$, and these BRs improve those that we gave in Ref. [7] (which were purely tree–level). Note that in Model I the BRs are independent of $\tan \beta$, while in Model I' there is a $\tan \beta$ dependence which causes the inequalities in Table 3 (we take $\tan \beta \geq 1.8$). For the MHDM it is not possible to predict the BRs since the parameters $X$, $Y$, and $Z$ may be varied independently of each other. The $cb$ channel is of order one percent in Models I and I' due to heavy CKM matrix suppression, although in a MHDM it is possible to enhance this channel ([6] and Section 7).

For all the charged scalars that we consider there exists an experimental lower bound from LEP of 54.5 GeV [23] which assumes $cs$ and $\tau\nu_\tau$ decays of $H^\pm$. This limit may not be valid if three–body decays dominate, although there exists a decay mode independent bound $M_{H^\pm} \geq 40$ GeV from considering visible decays of new particles.
Table 3: The branching ratios of $H^\pm$ excluding 3-body decays.

|                     | cs     | cb     | $\tau\nu_\tau$ |
|---------------------|--------|--------|-----------------|
| $H^\pm$ (Model I)   | 34.03% | 1.22%  | 64.75%          |
| $H^\pm$ (Model I')  | $\leq$ 4.76% | $\leq$ 0.17% | $\geq$ 95.07% |
| $H^\pm$ (MHDM)      | 0% $\rightarrow$ 100% | 0% $\rightarrow$ 100% | 0% $\rightarrow$ 100% |

contributing to the $Z$ width [24]. For the neutral sector there is no limit on $M_A$ since its standard production mechanism is in association with $h$, and so if $M_h$ is sufficiently large this channel would not be open. The literature mentions other ways of producing a light $A$ which we shall briefly review in Section 4. For $h$ the process $e^+e^- \rightarrow Z^{(*)} \rightarrow hf\bar{f}$ is available although this production method is proportional to $\sin^2(\beta - \alpha)$ and so may be suppressed. Hence for small values of $\sin^2(\beta - \alpha)$ a light $h$ with mass significantly lower than that of the current SM bound ($M_\phi \geq 87.6$ GeV) [3] is not ruled out. We note that in the MSSM it is possible to put actual limits of $M_h \geq 70.7$ GeV (all $\tan\beta$) and $M_A \geq 71.0$ GeV ($\tan\beta \geq 1$) [25]. For the heavier CP-even scalar ($H$) the presence of a small $\sin^2(\beta - \alpha)$ would automatically force $\cos^2(\beta - \alpha) \rightarrow 1$, thus enabling it to be produced with almost $\phi^0$ strength rate. However, in this scenario the decay $H \rightarrow hh$ may be open in all models and could dominate the standard $f\bar{f}$ decays. Assuming production via $e^+e^- \rightarrow Z^* \rightarrow HZ$, the signature would be a final state of 6 fermions, with four of them likely to be $b$ quarks coming from $h \rightarrow b\bar{b}$ decay. Such an event signature would be very similar to that coming from the process $e^+e^- \rightarrow Z h$ with subsequent decay $h \rightarrow AA$, which is considered in the searches, and so the current analysis should still be applicable to $H$.

We do not believe that $\rho$ parameter constraints in these scenarios of light Higgs bosons have been considered, especially bearing in mind the current bound on $M_{H^\pm} \geq 330$ GeV for the charged Higgs of Model II. One defines $\rho^0$ as:

$$\rho^0 = \frac{M_W^2}{\rho M_Z^2 \cos^2 \theta_W}$$

Here $\rho$ in the denominator contains all purely SM radiative corrections, while $\rho^0 \equiv 1$ in the absence of new physics. In the 2HDM there are extra contributions to $\rho^0$ [26] given by:

$$\Delta \rho^0 = \frac{G_F}{8\pi^2\sqrt{2}} \left[ \sin^2(\alpha - \beta) F(M_{H^\pm}^2, M_A^2, M_H^2) + \cos^2(\alpha - \beta) F(M_{H^\pm}^2, M_A^2, M_h^2) \right],$$

with

$$F(a, b, c) = a - \frac{bc}{b-c} \ln \frac{b}{c} - \frac{ab}{a-b} \ln \frac{a}{b} - \frac{ac}{a-c} \ln \frac{a}{c}.$$  

Ref. [27] shows that $-0.0017 \leq \Delta \rho^0 \leq 0.0027$ at the 2σ level. For the case of a light $h$ one requires $\sin^2(\beta - \alpha) \rightarrow 0$, and so the dependence on the heavier neutral CP-even
scalar $(H)$ drops out. If we demand a light $A$ no condition on $\sin^2(\beta - \alpha)$ is required, and so $M_H$ cannot be neglected from Eq. (5). Therefore the case of a light $h$ allows the above formula to simplify considerably and in Fig. 1 we plot the contribution $\Delta \rho^0$ against $M_A$ for three values of $M_{H^\pm}$, fixing $M_h = 20$ GeV. We see from Fig. 1 that to maintain $\Delta \rho^0$ within the $2\sigma$ limits (denoted by the horizontal lines) one requires $M_A \geq 250$ GeV in Model II, assuming $M_{H^\pm} \geq 330$ GeV. In Models I and I’ since $H^\pm$ may be light one may simultaneously allow a light $A$ and $h$ without violating the $2\sigma$ limits.

![Figure 1: $\Delta \rho^0$ as a function of $M_A$ with $M_h = 20$ GeV.](image)

4 A light $h$ or $A$ at LEP2

In this section we briefly review the current status of the literature on a light $A$ or $h$ i.e. the case of $M_h$ or $M_A \leq 40$ GeV. It is usually assumed that the sum of $M_A$ and $M_h$ is greater than $M_Z$, since the excluded region in the $M_A, M_h$ plane has the form $M_A + M_h \geq 90 \rightarrow 110$ GeV [25]. We note that the searches do not

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2On completion of this work we became aware of the new bound $M_{H^\pm} \geq 165$ GeV [28], [29]. This would eliminate the requirement $M_A \geq 250$ GeV.
consider the possibility of three-body decays to a lighter Higgs, although if such decays were present with a large BR they would have given a similar 6 fermion signature to those already searched for at LEP (see Section 5.3). Therefore we shall assume that $M_A + M_h \geq 90 \rightarrow 110$ GeV.

A light $h$ or $A$ has really only been considered in the context of the 2HDM (Model II) and has received limited attention $^{[30]} \rightarrow [34]$. These papers do not consider the three-body decay of a heavier Higgs boson to a light $h$ or $A$ plus a virtual vector boson, and this would be an alternative way of producing a $h$ or $A$ which has thus far escaped detection. We shall give particular attention to the 2HDM Model I and I', whose Higgs bosons may possess a significantly different phenomenology to those of Model II. In the following paragraphs we briefly summarize the existing methods of producing a light $h$ and $A$, and check to see if they are relevant for Models I and I'.

Ref. $^{[32]}$ studied the Yukawa production method, $e^+e^- \rightarrow b\bar{b}h(A)$. This process may be important in Model II in the case of large $\tan \beta$, since the $h(A)b\bar{b}$ coupling is proportional to $\tan \beta$. In Model I' the $h(A)b\bar{b}$ coupling is proportional to $\cot \beta$ and so one would have to rely on $h(A)$ emission from $\tau\tau$ pair production. This would give a much lower rate since the coupling $h(A)\tau\tau$ in Model I' is smaller than $h(A)b\bar{b}$ in Model II by a factor $3m^2_{\tau}/m^2_{b}$. In Model I all the fermion couplings scale as $\cot \beta$ and the Yukawa method would not be effective.

Ref. $^{[30]}$ considers the process $e^+e^- \rightarrow Z \rightarrow hhf\bar{f}$, whose dominant contribution comes from production of either $H$ or $H^*$ with subsequent decay to $hh$. If $H$ is on-shell then the rate can be large and is therefore a process in the same spirit to the methods we consider here – that is, producing $h(A)$ by the decay of an on-shell Higgs particle. This method was only considered for Model II, although it may also be used for Model I and I'. The $Hhh$ coupling is model independent, since this coupling originates from the Higgs potential. In the case of $H$ being on-shell, the BR($H \rightarrow hh$) will be different depending on the model in question. Applying this method to the case of a light fermiophobic Higgs (i.e. Model I with $\cos \alpha \rightarrow \pi/2$ and $\sin^2(\beta - \alpha)$ small) would cause the fermion couplings of $H$ to be scaled by a factor $1/\sin \beta$ relative to the SM Higgs. The production channel $e^+e^- \rightarrow Z^{(*)} \rightarrow HZ$ would then proceed with a rate close to that of the SM Higgs, and the subsequent decay $H \rightarrow H_FH_F$ would give the signature $\gamma\gamma\gamma\gamma$ and $f\bar{f}$ in opposite hemispheres. This topology would pass the current selection criteria for a $H_F$ which demands 2 isolated $\gamma$ recoiling against a fermion pair $^{[18]}$. In the case of $H$ being off-shell we still have a model dependent rate since the width of $H$ will appear explicitly in the propagator.

Refs. $^{[33]}$ and $^{[34]}$ consider the production of $AAA$ via $e^+e^- \rightarrow Ah^* \rightarrow AAA$ in the case when an on-shell $h$ cannot be produced. This process will again be model dependent for the reasons cited above. The case when the $h$ is on-shell is considered in the searches in Ref. $^{[25]}$, and we shall be covering this decay in more detail in Section 5.3.
5 Three–Body decays of Higgs Bosons

In the following sections 5.1 → 5.5 we study the impact of the three–body decays on the BRs and searches for $H^\pm$, $h$, and $A$ in the context of the 2HDM. We shall see that their importance varies from model to model and can be especially significant in Model I since $H^\pm$ and $A$ can decouple from the fermions at large $\tan \beta$ (see Table 1). We stress that an important use of these 3–body decays is that they may enable the detection of a light Higgs particle which has eluded current searches at LEP1 due to suppressed couplings to the $Z$.

5.1 The decays $H^\pm \to W^* h$ and $H^\pm \to W^* A$

In this subsection we consider the three–body decay $H^\pm \to h(A)f'\bar{f}$ mediated by a virtual $W^\pm$. Such a decay is not possible in the 2HDM (Model II) at LEP2 due to the bound $M_{H^\pm} \geq 330$ GeV. In Model I and $I'$ one may avoid the mass constraints from $b \to s\gamma$, as explained in Section 3, and so these models may possess a charged Higgs in the discovery range of LEP2. Studying the BR of this channel is important for two reasons:

(i) It may vastly alter current charged Higgs searches at LEP2, which always assume decays to $\tau\nu_\tau$ and $cs$.

(ii) It would be an alternative way of discovering a light $h$ or $A$ which is escaping current searches due to weak couplings to $Z$.

If the three–body decay channel for $H^\pm$ were dominant it would invalidate the current limit $M_{H^\pm} \geq 54.5$ GeV [23], and justify the use of the weaker limit $M_{H^\pm} \geq 40$ GeV [24]. Point (ii) is of particular interest for a fermiophobic Higgs which is searched for actively at the Tevatron and LEP, and may be hidden due to suppressed couplings to vector bosons. LEP [18] uses the standard Bjorken process, while the Tevatron [17] uses $q'\bar{q} \to H_F W$, which also depends on the vector boson coupling. In contrast, the cross–section $e^+e^- \to H^+H^-$ does not suffer mixing angle suppressions [33], instead being dictated by $M_{H^\pm}$. If $\text{BR}(H^\pm \to W^* h, (A))$ were large then it would be a copious source of a light $H_F$. We stress that this detection channel is not possible at LEP2 for a light $h$ and $A$ of the 2HDM (Model II) considered by Refs. [30] → [34] since $M_{H^\pm} \geq 330$ GeV. The partial widths for this channel are as follows [13], [36]:

$$\Gamma(H^\pm \to hW^\pm* \to hf\bar{f}) = \frac{9G_F^2 M_W^4}{16\pi^3} \cos^2(\beta - \alpha) M_{H^\pm} G_{hW}.$$ \hfill (8)

$$\Gamma(H^\pm \to AW^\pm* \to Af\bar{f}) = \frac{9G_F^2 M_W^4}{16\pi^3} M_{H^\pm} G_{AW}.$$ \hfill (9)
Figure 2: BR($H^± \to W^*h(A)$) in Model I against tan $\beta$ for $M_{h(A)} = 10$ GeV.

The functions $G_{AW}$ and $G_{hW}$ depend on the masses of the particles in the decay and are given in Refs. [19] and [37]. We note that the function $G_{ij}$ displayed in Refs. [19] and [37] (where $i$ and $j$ refer to a Higgs and a vector boson) contains a typing error, which is corrected by changing the final term $-2\lambda_{ij}/\kappa_j$ to $+2\lambda_{ij}/\kappa_j$. The function $G_{ij}$ is an approximation of the numerical integration over the Dalitz plot, but breaks down in the parameter space of interest to us ($M_{H^±} \leq M_W$). Therefore in the analysis that follows we shall evaluate $G_{ij}$ numerically. In the case of Eq. (8) the condition for a light $h$ causes $\cos^2(\beta - \alpha) \to 1$ and so enhances this width. Since we are interested in the case of a light $h$ we shall take $\cos^2(\beta - \alpha) \approx 1$, and so the results for the decay to $h$ and $A$ are more or less identical. We now consider in turn Models I and I'.

5.1.1 Model I

Fig. 2 displays BR($H^± \to W^*h(A)$) in Model I as a function of tan $\beta$ for 4 different values of $M_{H^±}$ (40 $\to$ 80 GeV), with $M_h(M_A) = 10$ GeV. As one can see, the BR of the three-body decay is close to 100% over the majority of the tan $\beta$ parameter space. The difference in mass between $M_{H^±}$ and $M_h(M_A)$ is important since it determines how off-shell the vector boson is, and so the curves for lower $M_{H^±}$ take longer to
reach $\approx 100\%$. In Fig. 3 we take $M_h(M_A) = 40$ GeV. We conclude that the three-body decays, if open, can be of great importance in Model I since the standard decays to 2 fermions are proportional to $\cot^2 \beta$. In addition, the possible large BR of this channel may have allowed $H^\pm$ to have avoided previous searches, and thus the limit $M_{H^\pm} \geq 40$ GeV should only be applied. The principal decays of $h$ and $A$ would be to $b\bar{b}$, unless $h$ is fermiophobic, in which case $H_F \to \gamma\gamma$ would dominate. In the case of fermiophobia one could find $\gamma\gamma f\bar{f}$ events in each hemisphere. LEP currently searches for $\gamma\gamma$ recoiling against $f\bar{f}$ \cite{18}, demanding isolated photons. Although in our new signature each photon pair would be accompanied by a pair of quarks, they would still pass the current selection criteria – in fact the efficiency of detecting any two of the four photons would increase by a factor $\approx 1.2$ \cite{18} relative to the efficiency for the $e^+e^- \to H_F Z \to \gamma\gamma f\bar{f}$ channel. We conclude that a very light $H_F$ could be copiously produced in $H^\pm$ decays and it should be possible from the topology to see whether one has registered a signal in the $e^+e^- \to H_F Z$ channel or $H_F H_F W^*W^*$ channel. The best discriminator would be the detection of 3 or more of the photons, which has a very small SM background and would have an efficiency $\approx 1.05$ times that for the $e^+e^- \to H_F Z$ channel. We stress that a light $H_F$ may continue avoiding searches in the $e^+e^- \to H_F Z$ channel due to weak coupling to the $Z$, although could be detected
via charged Higgs decays over a wide range of tan $\beta$ values, provided there are enough pair produced charged Higgs bosons. Lack of signal could be used to rule out regions of parameter space for $M_F/M_{H^\pm}/\tan \beta$.

As we pointed out in Ref. [39] an additional condition on the existence of a light $H_F$ is $M_F + M_A \geq 160$ GeV. This is due to the fact that the channel $e^+ e^- \rightarrow AH_F \rightarrow \gamma\gamma f\bar{f}$ is also searched for in Ref. [18] and is complementary to $e^+ e^- \rightarrow H_F Z \rightarrow \gamma\gamma f\bar{f}$. Therefore the former must be closed kinematically if one wishes to consider suppressed $e^+ e^- \rightarrow H_F Z$ production and the possibility of a light $H_F$.

If $h$ is not fermiophobic then one could have final states of 8 fermions, with four of them likely to be $b$ quarks. Such a signature might allow detection of $H^\pm$ in the difficult $M_{H^\pm} \approx M_W$ region, although a full analysis is beyond the scope of this paper.

5.1.2 Model I'

In Model I' the three–body decay can only be strong only at small tan $\beta$, which we can see from Fig. 4. This is because the $H^\pm \rightarrow \tau \nu\tau$ decay width is proportional to $\tan^2 \beta$ and so $H^\pm$ does not decouple from the fermions as tan $\beta$ increases. Regarding the possible signatures, $h(A) \rightarrow \tau\tau$ and $h(A) \rightarrow b\bar{b}$ would have equal rates for tan $\beta \approx 2$, with the former quickly approaching 100% as tan $\beta$ increases. Hence one would expect a final state of 8 fermions, with 4 of them likely to be $\tau$ leptons. Note that we are interested in the region $\alpha \approx \beta$ in order to ensure a light $h$, and so the $h \rightarrow \tau\tau, b\bar{b}$ couplings lose their $\alpha$ dependence and depend purely on tan $\beta$.

5.2 The decay $h \rightarrow H^\pm W^*$

For the decay $h \rightarrow H^\pm W^*$ one can only consider Models I and I', since in Models II and II' the bound $M_{H^\pm} \geq 330$ GeV disallows this process at LEP2. The production mechanisms for $h$ are $e^+ e^- \rightarrow Z^* \rightarrow hZ$ and (if open) $e^+ e^- \rightarrow Z^* \rightarrow Ah$. These two channels are complementary, the former being proportional to $\sin^2(\beta - \alpha)$ and the latter to $\cos^2(\beta - \alpha)$.

The couplings of $h$ to $f\bar{f}$ involve the mixing angle $\alpha$ (see Table 1), and in Model I have a minor dependence on tan $\beta$; this is because the factor $(\sin \beta)^{-1}$ appears explicitly and for tan $\beta \geq 2$ takes values between 1 and 1.12. In Fig. 5 (for Model I) we plot BR($h \rightarrow H^\pm W^*$) as a function of cos $\alpha$, fixing tan $\beta = 2$, for values of $M_h$ up to 100 GeV (we take $M_{H^\pm} = 55$ GeV). From the figure we see that the BR for the three–body decay is at a maximum at $\alpha = \pi/2$, (the condition for fermiophobia), and for low values of cos $\alpha$ the competing decays are $h \rightarrow \gamma\gamma$ and $h \rightarrow WW^*$. As cos $\alpha$ increases the standard decays $h \rightarrow f\bar{f}$ start to gain in strength.

In Model I' there is a strong tan $\beta$ dependence, since the coupling $h \rightarrow \tau\tau$ grows as tan $\beta$. We find that BR($h \rightarrow H^\pm W^*$) peaks at $\approx 5\% (\approx 0.5\%)$ for $M_h = 100$ GeV with
5.3 The decay $h \rightarrow Z^* A$

We now consider the case of $M_A \leq M_h$ which allows the decay $h \rightarrow AZ^*$, and we shall assume that decays to $H^\pm W^*$ are not open. If $M_h \geq 2M_A$ the channel $h \rightarrow AA$ would be open as well. The experimental signature of these two decays would be similar, both giving 6 fermion final states coming from either production mechanism. There are four different topologies here which we list in Table 4, and the colon separates the particles into hemispheres. We note that in principle these 4 distinct topologies could be distinguished by measuring the invariant masses of the jets and/or using $b$-tagging.

| $e^+e^- \rightarrow hZ$ | $Z : Aff$ | $Z : AA$ |
|------------------------|----------|----------|
| $e^+e^- \rightarrow hA$ | $Aff : A$ | $A : AA$ |

Table 4: The 6 fermion topologies originating from $h$ and/or $A$ production.
techniques. The current searches do consider 6 fermion final states, but from the lack of signal rule out regions in the $M_h - M_A$ plane only for $M_h \geq 2M_A$. For the region $M_h \leq 2M_A$ only the decays $h \rightarrow far{f}$ are considered. We wish to see if the three-body decays (and thus the 6 fermion signature) can be important even when the decay $h \rightarrow AA$ is not open. In the following subsections we consider the strength of the decay $h \rightarrow AZ^*$ in Models I, I' and II respectively.

5.3.1 Model I

In Fig. 6 (again for $\tan \beta = 2$) we plot $\text{BR}(h \rightarrow AZ^*)$ in Model I for values of $M_A$ which do not permit $h \rightarrow AA$ (we take $M_h = 2M_A - 5$ GeV). If $h \rightarrow AA$ is open, it will contribute to the 6 fermion signature although giving a different topology as shown in Table 4. Note that the production process $e^+e^- \rightarrow Z^* \rightarrow hZ$ would not be open for the values of $M_h \geq 100$ GeV displayed in Table 4. We see from Fig. 6 that lighter values of $M_h$ allow larger $\text{BR}(h \rightarrow AZ^*)$ at small values of $\cos \alpha$ than for heavier $M_h$; as $\cos \alpha$ increases the curves for lighter $M_h$ fall more rapidly. This can be explained as follows. Larger $M_h$ increases the partial width for the three-body decay (since the $Z$ will be less off-shell), but also enhances the decay $h \rightarrow WW^*$. Hence in the small
Figure 6: BR($h \to AZ^*$) against $\cos \alpha$ for Model I when $h \to AA$ is not allowed.

In the $\cos \alpha$ region the curves with larger $M_h$ take lower BR values since in this region the competing decay is $h \to WW^*$, with the fermions decoupled. As $\cos \alpha$ increases the fermion decays start to gain in importance, and so all the curves fall to small values; those curves with larger $M_h$ do not fall as sharply since their partial width is larger. Of course, one may choose $M_h$ and $M_A$ such that the on–shell decay $h \to ZA$ is open.

For these choices of Higgs masses the decay $h \to AA$ would always be open (since $2M_A \leq M_Z + M_A$) and so one would expect the 6 fermion signature to dominate the 4 fermion signature.

5.3.2 Model I’ and Model II

In both these models the fermions never decouple completely from $h$, and so we expect lower maximum BRs for the three–body decay. In addition there will be a strong $\tan \beta$ dependence. In Model I’ we find a peak BR($h \to AZ^*$) $\approx 10\%$ $(1\%)$ for $\tan \beta = 2(10)$, with $M_h = 120$ GeV. In Model II these values drop to $\approx 5\%$ and $\approx 0.2\%$ respectively.
5.4 The decay $A \rightarrow Z^*h$

These decays have not been considered so far in the searches for $A$ of the general 2HDM, and if open may be of great importance in Model I. The phenomenology of $A$ in Model I and Model II has been considered in the context of the LHC [10], but this work did not consider three–body decays, and was concerned with a more massive $A$ which could decay to on–shell $hZ$. The condition for a light $h$ causes $\cos^2(\beta - \alpha) \rightarrow 1$,

![Figure 7: BR($A \rightarrow Z^*h$) in Model I against $\tan \beta$ for $M_h = 40$ GeV.](image)

and so in this case the production cross-section $e^+e^- \rightarrow Z^* \rightarrow Ah$ would be not mixing angle suppressed. We now consider the strength of the decay ($A \rightarrow Z^*h$) in Models I, I' and II.

5.4.1 Model I

In Fig. 7 we plot BR($A \rightarrow Z^*h$) for 3 values of $M_A$, fixing $M_h = 40$ GeV. We require $M_A + M_h \geq 110$ GeV (from LEP1 searches), and deliberately consider mass choices which make $Z$ off–shell in order to emphasize that the off–shell decays can be prominent. We stress that lower values of $M_h$ would imply that the on–shell decay $A \rightarrow hZ$ is open and so the three–body BR would be $\approx 100\%$ over all the $\tan \beta$ range.
– for this reason we do not plot a graph. We also assume in Fig. 7 that $M_{H^\pm}$ is not light enough to cause a competing three–body decay. From Fig. 7 one can see the importance of the decay $A \rightarrow Z^* h$ when open, and it would be an alternative way of producing a light $H_F$ with a good rate. The signature would be similar to that discussed in Section 5.1.1, although in this case there would be less jets.

Comparing Fig. 7 with Fig. 3 one sees a significant difference in the strength of the respective lines for $M_{H^\pm} = 80$ GeV and $M_A = 80$ GeV, the former starting at $BR \approx 50\%$ and the latter at $BR \approx 5\%$. There are two reasons why the three–body decay for $H^\pm$ is stronger than that for $A$, for identical masses of the Higgs bosons in the decay, despite the fact that the $Z$ mediated decay has a slightly stronger coupling by a factor of $\approx 4/3$. Firstly, the fact that $M_Z \geq M_W$ means that $Z$ would be more off–shell in the decay $A \rightarrow hZ^*$ than $W$ in the decay $H^\pm \rightarrow W^* h (A)$. Secondly, and more importantly, $BR(H^\pm \rightarrow cb)$ is strongly suppressed by the CKM matrix while the decay $A \rightarrow b\bar{b}$ is normally the dominant decay for $A$. 
5.4.2 Model I' and II

In Fig. 8 we plot the analogy of Fig. 7 for Model I'. We find a similar tan β dependence to that found for $H^\pm \rightarrow W^* h(A)$ in Fig. 4, with a decrease in the BR as tan β increases. The curve with $M_A = 140$ GeV allows the two–body decay $A \rightarrow hZ$, and maintains a large BR until tan β becomes very large.

In Model II one may consider the decay $A \rightarrow hZ^*$, although since $M_{H^\pm} \geq 330$ GeV in this model the ρ parameter constraints (see Fig. 1) suggest that requiring a light $h$ would cause $M_A$ to be out of the range of LEP2. For this reason we do not plot a graph. Even if one allowed $M_A$ in range at LEP2 the three–body decays would be very small, this being due to the fact that the decay width for $A \rightarrow b\bar{b}$ is proportional to tan$^2$ β in Model II. If an on–shell $Z$ is allowed (e.g. $M_A = 140$ GeV, $M_h = 40$ GeV) one finds BR($A \rightarrow hZ$) ≈ 70% at tan β = 2, falling below 1% at tan β = 30.

![Graph](image)

Figure 9: BR($A \rightarrow H^\pm W^*$) in Model I against tan β, for $M_H = 55$ GeV.

5.5 The decay $A \rightarrow H^\pm W^*$

The decay $A \rightarrow H^\pm W^*$ may be considered in Models I and I'. In Model II the bound $M_{H^\pm} \geq 330$ GeV disallows this channel for $M_A$ in range at LEP2. In Fig. 9 (for
Model I) we vary $M_A$ from $70 \rightarrow 100$ GeV, setting $M_{H^\pm} = 55$ GeV. We do not consider values of $M_A \geq 100$ GeV, since we require that the production process $e^+e^- \rightarrow hA$ be open, and also that $M_h \geq M_A$ so that there are no competing $A \rightarrow hZ^*$ decays. We see that this channel can be dominant at larger values of $\tan \beta$. For Model I' with $M_A = 100$ GeV the BR peaks at $\approx 5\%$, and falls to below $1\%$ for $\tan \beta \geq 10$ (we do not plot a graph).

6 Detection in an enhanced $H^\pm cb$ channel

We now consider the $H^\pm$ of the MHDM and aim to give a more detailed analysis of a detection channel mentioned in our earlier work. If no neutral Higgs boson exists with a mass lighter than $M_{H^\pm}$ the BRs of $H^\pm$ in Model I and I' will be given by the entries in Table 3. In the MHDM there is a possibility of an enhanced $cb$ decay, which was first mentioned in Ref. [5] and in our earlier work [7] we displayed the parameter space which allowed large $\text{BR}(H^\pm \rightarrow cb)$. We note that we did not use running masses of the quarks in Ref. [7], and Figures 1 and 2 there are for $m_c = 1.5$ GeV and $m_b = 5$ GeV. At the energy scale of 100 GeV these values drop to $m_c = 0.62$ GeV and $m_b = 3.04$ GeV, and so our previous results slightly overestimate $\text{BR}(H^\pm \rightarrow cb)$. However, it will still be possible to have a large $\text{BR}(H^\pm \rightarrow cb) \geq 10\%$ if $|X| \geq |Y|, |Z|$ by a factor of 5 or so.

A review of the detection techniques for a light $H^\pm$ at LEP2 appears in Ref. [11]. The three signatures analysed were $cscs$, $cst\nu_r$ and $\tau\nu_r\tau\nu_r$. In the MHDM one could consider the signature $cb\tau\nu_r$, which would give an isolated $\tau$ and missing energy recoiling against a hadronic system with a tagged $b$ quark. Much of the selection criteria would be identical to that of the $cst\nu_r$ channel. In the $cst\nu_r$ channel the $W^+W^-$ background is removed by reconstructing the invariant masses of the $cs$ and $\tau\nu_r$ systems, thus significantly reducing this background if $M_{H^\pm} \leq M_W$. The expected event numbers/efficiencies for the $cst\nu_r$ channel are given in Table 5 (from Ref. [11]).

| Process | Eff. or No. of backg. evts |
|---------|---------------------------|
| $H^\pm$ (60 GeV) | 5.6% |
| $q\bar{q}$ | 0 |
| $W^+W^-$ | 2 |
| $ZZ$ | 0 |
| $\tau^+\tau^-$ | 0 |

Table 5: Expected signal efficiency and background event numbers in the $cst\nu_r$ channel for $\sqrt{s} = 175$ GeV.

Detection of a $H^\pm$ will be very difficult in the region $M_H \approx M_W$, partly due to the lack of $H^+H^-$ pair production events for this mass region, and partly due to the fact
that invariant mass cuts which reduce the $WW$ background also remove the signal. In Ref. [7] we showed that the presence of a large BR($H^\pm \to cb$) (of order 50% is possible) in a MHDM has two potential uses:

(i) It would suggest that a detected $H^\pm$ is from the MHDM, since the equivalent decay in Model I and Model I’ has a BR less than 1%.

(ii) It may allow a chance detection in the difficult $M_H \approx M_W$ region since $W^\pm$ rarely decays to $cb$.

The $c\tau\nu_\tau$ channel has the advantage of almost negligible background from $WW$ and $ZZ$ events, as long as one can correctly identify the $b$ quark. In addition we will be able to use the selection criteria of the $cs\tau\nu_\tau$ channel, replacing the invariant mass cut with a $b$–tag requirement. The number of signal events (before cuts) in the $c\tau\nu_\tau$ channel is given by:

$$N_{H^+H^-} \times f(BR),$$

where $N_{H^+H^-}$ is the number of pair produced $H^\pm$ events, and $f(BR)$ is defined by

$$f(BR) = 2 \times \text{BR}(H^\pm \to cb) \times \text{BR}(H^\pm \to \tau\nu_\tau).$$

In order to isolate these final states $b$–tagging will be necessary. Since this a standard technique for searching for the SM Higgs at LEP2, the efficiency, $e_b$, will be quite high ($\approx 70\%$) in practice, see for example Ref. [12]. For the $c\tau\nu_\tau$ channel we shall use the optimistic values $\text{BR}(H^\pm \to cb) = \text{BR}(H^\pm \to \tau\nu_\tau) = 50\%$ which maximizes $f(BR)$. We note that in the $cs\tau\nu_\tau$ channel the invariant mass cut was the last cut applied, and all other non–$WW$ backgrounds had already been removed. When it is applied for the mass region $M_{H^\pm} \approx M_W$ it will remove the signal as well. Our aim is to replace this latter cut with a $b$–tag requirement, which will reduce the $WW$ background to negligible proportions while preserving most of the signal. One would need to have a strong rejection against fake tags coming from $c$ quarks. From the figures in Ref. [11] one can infer that the invariant mass cut reduces the Higgs signal by $\approx 2/3$. Therefore we shall assume a selection efficiency ($e_H$) before $b$–tagging of 8.4% for the Higgs signal, obtained by scaling the value of 5.6% in Table 5. One can obtain the following formula for the number of signal events ($N_{\text{sig}}$) in the $c\tau\nu_\tau$ channel:

$$N_{H^+H^-} \times e_H \times e_b \times f(BR).$$

The number of $c\tau\nu_\tau$ events from $WW$ production is 1.1 before any cuts have been applied. The cuts before the $b$–tag requirement have a selection efficiency considerably below 100% and therefore the background is entirely negligible. We stress that we require strong $c$ quark rejection since the number of events in the $cs\tau\nu_\tau$ channel is large. We then require 3 or more signal events for detection, and we see from the
Eq. (12) that the number of events for \( M_{H^\pm} = 80 \text{ GeV} \) and \( \sqrt{s} = 180 \text{ GeV} \) is equal to 2.1. Therefore detection is certainly marginal. At the higher collider energy, \( \sqrt{s} = 200 \text{ GeV} \), due to larger \( N_{H^+H^-} \) we find that \( N_{\text{sig}} = 3.8 \). All this analysis is with optimistic choices for \( e_b \) and \( f(BR) \). With greater luminosity, which would be available at a next generation collider one could probe a greater parameter space of \( f(BR) \).

In summary, the \( cb\tau\nu_\tau \) channel at LEP2 only provides a slight chance of overcoming the difficult \( M_W \approx M_{H^\pm} \) region in the MHDM, since the largest values of \( f(BR) \) would be needed. However, the signature would have a use for \( M_{H^\pm} \) lighter than \( M_W \) since it would provide evidence of the MHDM. For \( M_{H^\pm} \) comfortably below \( M_W \) a reasonable number of \( H^\pm \) pairs would be produced and we shall require three tagged \( cb\tau\nu_\tau \) events to conclude that a detected \( H^\pm \) originates from the MHDM. From Eq. (12) we can obtain Eq. (13) from which the values of \( f(BR) \) needed to produce the distinctive signature of three tagged \( cb\tau\nu_\tau \) events can be found.

\[
f(BR) \times N_{H^+H^-} = 36.
\]

Thus for \( N_{H^+H^-} = 100 \) (corresponding to \( M_H = 75 \text{ GeV} \) at \( \sqrt{s} = 180 \text{ GeV} \)), one finds \( f(BR) = 0.36 \), and so \( \text{BR}(H^\pm \rightarrow cb) \geq 20\% \) is required. For lower masses \( \text{BR}(H^\pm \rightarrow cb) \approx 10\% \) (or even less) would be sufficient. We note that for the 2HDM that the analogous signal would be \( \leq 0.1 \) events, which is unobservable.

Another way of distinguishing \( H^\pm \) of the MHDM would be through a lack of \( H^\pm \rightarrow \tau\nu_\tau \) decays i.e. leptophobia with \( \text{BR}(H^\pm \rightarrow jets \rightarrow 100\%) \). From Table 1 one sees that \( \text{BR}(H^\pm \rightarrow \tau\nu_\tau) \) is expected to be large in both Model I and I' (\( \approx 65\% \) and \( \geq 95\% \) respectively), while in the MHDM it can be reduced to much lower values. For example, a simple calculation shows (not including \( cb \) decays)

\[
\text{BR}(H^\pm \rightarrow \tau\nu_\tau) \approx \frac{1.8|Z|^2}{1.8|Z|^2 + |Y|^2}.
\]

Therefore if \( |Y| \geq 2|Z| \), one finds \( \text{BR}(H^\pm \rightarrow \tau\nu_\tau) \leq 30\% \). Including the \( cb \) decays would reduce this further, and so it is apparent that a sizeable parameter space exists for \( \text{BR}(H^\pm \rightarrow \tau\nu_\tau) \ll 65\% \). For the extreme case of \( \text{BR}(H^\pm \rightarrow jets) \rightarrow 100\% \) one would find a ninefold increase in the number of events in the \( cs\bar{c}s \) channel compared to Model I.

7 Conclusions

We have studied the impact of three–body decays of a Higgs boson to a lighter Higgs boson and a virtual vector boson in the context of the non–supersymmetric 2HDM model. Such decays have been studied in the MSSM, although their importance is
magnified in the four versions of the general 2HDM, partly due to the lack of correlation among both the mixing angles and masses of the Higgs scalars, and partly due to the different couplings of the Higgs bosons to the fermions. Such decays would allow the production of light neutral Higgs bosons which have so far escaped detection at LEP due to suppressed couplings to the Z boson. We showed that the three–body channels, if open, can be of great importance in Model I, permitting the decays $H^\pm \rightarrow h(A)W^*$ and $A \rightarrow hZ^*, H^\pm W^*$ to proceed with large branching ratios over a wide range of $\tan \beta$ values. As $\tan \beta$ increases $H^\pm$ and $A$ decouple from the fermions and so the three–body channels rapidly grow in importance, ultimately reaching branching ratios close to 100%. These results have important applications for the phenomenology of Model I at LEP2, particularly for $H^\pm$ which may avoid the present search techniques which assume charged scalar decays to $cs$ or $\tau\nu_\tau$. We suggested that the three–body decays might also allow detection of $H^\pm$ in the difficult $M_{H^\pm}$ region. A fermiophobic Higgs boson (only possible in Model I) which has so far escaped direct searches in the $e^+e^- \rightarrow H_FZ$ channel may be produced copiously in the decay of $H^\pm$ or $A$, provided that enough of the latter are produced on–shell. Signatures with 3 or 4 photons would be possible, which would pass current $H_F$ search criteria with equal or better efficiencies, and enable discrimination from the $e^+e^- \rightarrow H_FZ$ signal. For the CP–even $h$ the three–body decays $h \rightarrow H^\pm W^*$ and $h \rightarrow AZ^*$ may be dominant at low values of $\cos \alpha$ and allow a 6 fermion signature even when the decay $h \rightarrow AA$ is not open.

In Model I’ the three–body decays can be significant (although not usually dominant) for small values of $\tan \beta$. As $\tan \beta$ increases the decays of $H^\pm$, $A$ and $h$ to the third generation of leptons dominate. In Model II one finds smaller branching ratios than the analogous cases for Model I’, and the decays involving $H^\pm$ are not relevant at LEP2 energies due to the bound of $M_{H^\pm} \geq 330$ GeV. Finally, we showed that a $H^\pm$ in a general MHDM with $N$ doublets may be distinguished from $H^\pm$ of a 2HDM if it possessed a sizeable ($\geq 10\%$) $\text{BR}(H^\pm \rightarrow cb)$ or a $\text{BR}(H^\pm \rightarrow jets) \gg 30\%$.

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