LOWER BOUNDS ON CHARGED HIGGS BOSONS
FROM LEP AND THE TEVATRON

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ABSTRACT

We point out that charged Higgs bosons can decay into final states different from $\tau^+\nu_\tau$ and $c\bar{s}$, even when they are light enough to be produced at LEP II or at the Tevatron through top quark decays. These additional decay modes are overlooked in ongoing searches even though they alter the existing lower bounds on the mass of the charged Higgs bosons that are present in supersymmetric and two Higgs doublets models.
The discovery of a charged Higgs boson would be an unambiguous signal of an extended Higgs sector and possibly of supersymmetry. In supersymmetric models, at least two Higgs doublets are needed to give mass to all fermions: one is coupled only to down-type quarks and leptons; the other, only to up-type quarks. A Two Higgs Doublet Model (2HDM) is said of Type II if the doublets are coupled as in supersymmetric models with minimal particle content. It is said of Type I if one Higgs doublet does not couple to fermions at all and the other couples as the Standard Model (SM) doublet.

After electroweak symmetry breaking, five physical states remain: two CP-even Higgs bosons $h$ and $H$ (with $m_h < m_H$), a CP-odd Higgs boson $A$, and two charged states $H^\pm$. The charged-Higgses–fermions interactions, can then be comprehensively expressed as:

$$\mathcal{L} = \frac{g}{\sqrt{2}} \left\{ \left( \frac{m_{dL}}{M_W} \right) X \overline{u}_L V_{ij} \, d_{Ri} + \left( \frac{m_{uL}}{M_W} \right) Y \overline{u}_{Ri} \, V_{ij} \, d_{Lj} + \left( \frac{m_{lL}}{M_W} \right) Z \overline{\nu}_{Li} \, \nu_{Ri} \right\} H^+ + \text{h.c.}, \quad (1)$$

where $V$ is the CKM matrix. The equality $X = Z = 1/Y = \tan \beta$, with $\tan \beta$ the ratio of the two vacuum expectation values, identifies 2HDMs of Type II and supersymmetric models; $Y = -X = -Z = \cot \beta$, identifies 2HDMs of Type I.

Besides the mass of $h$, $H$, $A$, and $H^\pm$, two additional parameters are needed to describe the Higgs sector in 2HDMs of Types I and II: $\tan \beta$ and the mixing angle $\alpha$. In supersymmetric models, the Higgs sector is more constrained, and only two free parameters are needed at the tree level, $m_A$ and $\tan \beta$. Supersymmetry induces a relation between $\tan 2\beta$ and $\tan 2\alpha$ and the well-known tree-level sum rule $m_{H^\pm}^2 = m_A^2 + m_H^2$, which is only mildly altered by one-loop corrections [1]. Together with the experimental lower bound on $m_h$, $m_A > 92 \text{ GeV}$, for $\tan \beta > 1$ [2], this sum rule makes the supersymmetric charged Higgs bosons possible candidates for discovery at the Tevatron, but not at LEP II.

Strong constraints on charged Higgs bosons come from searches of processes where $H^\pm$ is exchanged as a virtual particle. Among them, the measurement of the inclusive decay $B \to X_s \gamma$ [3] excludes charged Higgs bosons in a 2HDM of Type II up to $\sim 165 \text{ GeV}$ [3]; however it is, in general, inconclusive for supersymmetric models [4] and 2HDMs of Type I [3, 4]. Other indirect bounds on the ratio $m_{H^\pm}/\tan \beta$ come from inclusive semileptonic $b$-quark decays $B \to D\tau\nu_\tau$, $m_{H^\pm} \gtrsim 2.2 \tan \beta \text{ GeV}$ [5] and from $\tau$-lepton decays, $m_{H^\pm} \gtrsim 1.5 \tan \beta \text{ GeV}$ [5]. They apply to charged Higgs bosons of Type II in 2HDMs and supersymmetric models. In the former, however, they are non-competitive with the stronger lower bound due to the measurement of $B \to X_s \gamma$; in the latter they are already saturated by the above sum rule and the lower bound on $m_A$. Constraints on the low-$\tan \beta$ region and light $H^\pm$ in Type I models come from the measurement of $Z \to b\bar{b}$ and $B^0-\bar{B}^0$ mixing (see discussion in [5]).

It is possible that the 2HDMs described above are only “effective” models, i.e. the low-energy remnant of Multi-Higgs-Doublets models, with the same number of degrees of physical states non-decoupled at the electroweak scale. In this case, more freedom remains in the possible values that $X$, $Y$, and $Z$ can acquire. For $X = -1/Y = -a$, with $a \geq 2$, for example, charged Higgs bosons with $m_{H^\pm} = 100 \text{ GeV}$ can escape the $\bar{B} \to X_s \gamma$ constraint [4], while having widths for decays into light fermions substantially coinciding with those obtained in a 2HDM of Type II. Moreover, lepton and quark couplings in (1)
may be unrelated, thus rendering the indirect bounds from b-quark and τ-lepton decays independent of that coming from $B \to X_s\gamma$. Indirect and direct bounds are, therefore, all equally necessary in providing the complementarity that allows the exclusion of certain ranges of $m_{H^\pm}$ in supersymmetric models, in Type I and Type II 2HDMs, and in those models that may counterfeit them in one specific search.

Charged Higgs bosons are searched for at LEP II, above the LEP I limit, in the range $45 \lesssim m_{H^\pm} \lesssim 100$ GeV and at the Tevatron in the range $m_{H^\pm} < m_t - m_b$, i.e. when produced by a decaying $t$-quark. Searches at LEP II rely on the assumption that no $H^+$ decay mode, other than $c\bar{s}$ and $\tau^+\nu_\tau$, is kinematically significant; they give a limit $m_{H^\pm} \gtrsim 78.6$ GeV [4], which applies to 2HDMs of Types II and I. Indeed, within the assumption $BR(H^+ \to c\bar{s},\tau^+\nu_\tau) \approx 100\%$, in Type I models the two branching ratios are tan$\beta$-independent and approximately equal to those obtained in Type II models with tan$\beta = 1$.

At the Tevatron, searches of an excess of $t\bar{t}$ events in the $\tau$ channel provide a tan$\beta$-$m_{H^\pm}$ exclusion contour that constrains the very-large-tan$\beta$ region in supersymmetric models and 2HDMs of Type II [10], for which the rate of $t \to H^+b$ is large. Similarly large is this rate in the region of low tan$\beta$ (tan$\beta \lesssim 1$), for Type II Yukawa couplings. Searches of $H^+$ apply in this region to the non-supersymmetric case. They are carried out, specifically for this type of couplings, looking for: i) a deficit in the $e,\mu$ channels, due to $H^+ \to c\bar{s}$, for $m_{H^\pm} \lesssim 130$ GeV; ii) a larger number of taggable $b$-quarks due to $H^+ \to t^*b \to b\bar{b}W$ for $m_{H^\pm} \gtrsim 130$ GeV [11, 12]. Given the limited luminosity at present available at the Tevatron ($\sim 1$ fb$^{-1}$), there is no sensitivity to the intermediate range of tan$\beta$ where the rate $t \to H^+b$ becomes low. This region, partially accessible at the upgraded Tevatron, will be fully covered at the LHC [13].

The aim of this letter is to show that there exist additional decay modes, which are overlooked in ongoing searches of $H^+$ within 2HDMs and supersymmetric models, and which alter the existing lower bounds on $m_{H^\pm}$. In the following, the considered type of weak scale supersymmetry has minimal particle content and R-parity conservation. No specific assumption is made on the superpartner spectrum and on the scale/type of messengers for supersymmetry breaking. All branching ratios presented for supersymmetric models are calculated using HDECAY [14].

In 2HDMs, these modes are $H^+ \to AW^+$ and/or $hW^+$ ($HW^+$). They produce mainly the same final state $b\bar{b}W^+$, as the above-mentioned $bt^*$ mode and, to a lesser extent, the state $\tau^+\tau^-W^+$. Our statement is based on the fact that there is no stringent lower bound on $m_A$ and/or $m_h$ coming from LEP [15]. Indeed, since the mixing angle $\alpha$ is, in this case, a free parameter, one can think of a scenario in which the coupling $ZhA$ vanishes. This coupling being proportional to $\cos(\beta - \alpha)$, the required direction is $\alpha = \beta \pm \pi/2$. In this case, the process $Z^* \to hA$ does not occur and the LEP II bound $m_A > 92$ GeV obtained for supersymmetric models does not hold. Nevertheless, the cross section for the process $e^+e^- \to Z^* \to hZ$, proportional to $\sin^2(\beta - \alpha)$, is not suppressed with respect to that for the corresponding production mechanism of the SM Higgs boson, and the LEP II bound $m_h > 114$ GeV [2] applies to our case. The coupling $ZH A$, still proportional to $\sin(\beta - \alpha)$, has also full strength, whereas $HZZ$ vanishes. The process $Z^* \to HA$ could
in principle provide a bound on $m_A$ depending on $m_H$ and $\tan \beta$. For large $m_H$, however, no real lower bound can be imposed on $m_A$. Conversely, even without making specific choices on the angle $\alpha$, one can assume $h$ to be heavy enough to render impossible any significant lower bound on $m_A$. The other two production mechanisms possible at LEP I (they require larger numbers of events than LEP II can provide) are the decay $Z \to A\gamma$ and the radiation out of $b\bar{b}$ and $\tau^+\tau^-$ pairs [16]. The first is mediated only by fermion loops, unlike the decay $Z \to h\gamma$, which has additional contributions from $W$-boson loops. The corresponding rate is about two orders of magnitude smaller than that for $Z \to h\gamma$ and therefore too small to allow for a visible signal [17]. The second process allows for sizeable rates only for very large values of $\tan \beta$. No bound can be obtained for non-extreme values of $\tan \beta$ and for 2HDMs of Type I. In general, therefore, one remains with the rather modest bound from the decay $\Upsilon \to A\gamma$, which has been searched for by the Crystal Ball Collaboration [18], $m_A > 5 \text{ GeV}$.

\begin{center}
\includegraphics[width=\textwidth]{figure1}
\end{center}

\textbf{Figure 1:} Branching fractions for the decay $H^\pm \to AW^*$ as a function of $m_A$ for three values of $m_{H^\pm} = 70, 110$ and 150 GeV and two values $\tan \beta = 1$ (solid) and 2 (dotted).

If one recalls that the interaction term $H^+W^-A$ is weighted by a gauge coupling, unsuppressed by any projection factor, it is clear that the decay $H^+ \to AW^+$ can be rather important for Type I models, or for models of Type II with small $\tan \beta$. This remains true even for an off-shell $W$-boson, in spite of the additional propagator and weak coupling that are then required. For a 2HDM of Type I and Type II with $\tan \beta = 1$, the branching ratios $BR(H^+ \to AW^+)$ are shown in Fig. 1 as functions of $m_A$ for different values of $m_{H^\pm}$ (solid lines). Already for $m_{H^\pm} = 70$ GeV, roughly the lower bound obtained at LEP II when $BR(H^+ \to c\bar{s}, \tau^+\nu_\tau) \simeq 100\%$ is assumed, the branching ratio is 50\%–20\% for $m_A = 10$–30 GeV. More strikingly, for heavier $H^\pm$, when the $W$-boson is not too far from being on shell, this decay mode becomes the dominant one. We also show in Fig. 1 the branching ratios for this decay mode in a Type II model with a higher value $\tan \beta = 2$ (for Type I model, the situation does not change). $BR(H^\pm \to W^*A)$ is of course smaller
because the competing decay mode, $H^+ \rightarrow \tau^- \nu$, has an enhanced decay width. This is more striking for low $m_{H^\pm}$ values when the $H^\pm \rightarrow AW^*$ decay channel occurs only at the three-body level. For a heavier $H^\pm$ boson, values of $\tan \beta$ slightly larger than unity do not change the main trend. This is particularly true when the $W$ boson is on–shell as in the example with $M_H \sim 150$ GeV and a light pseudoscalar $A$ boson. In this case, only for much larger $\tan \beta$ values that the $H^\pm \rightarrow AW^*$ decay mode becomes dominant and then, the search for the $H^\pm$ boson at LEP II will be the standard one and the limit $m_{H^\pm} \geq 78.6$ GeV form $\tau \nu$ and $cs$ decay \[9\] will hold (for intermediate $\tan \beta$ values, one has to take into account simultaneously all decay modes, rendering the analysis more complicated).

Since the two modes $hW^+$ and $HW^+$ are forbidden respectively by our choice of $\alpha$ and the requirement of a very heavy $H$, the other competing channels are $\tau^\nu$ and $\bar{c}s$ in the Tevatron searches. In Fig. 2, the final branching ratio $BR(H^\pm \rightarrow \bar{b}bW^*)$ is shown as a function of $m_{H^\pm}$ in a 2HDM of Type II, with our choice of $\alpha$, for different values of $\tan \beta$ and of $m_A$. For the larger $m_A$, the mode $AW^+$ is forbidden. Indeed, above $m_{H^\pm} = 130$ GeV the mode $\bar{c}s$ is quickly taken over by $\bar{b}t^*$, with the same $\tan \beta$ dependence, but much larger Yukawa couplings, which can compensate the virtuality of the $t$-quark. The deviations from this pattern become striking when the mode $AW^+$ starts being allowed.

![Figure 2: Branching fractions for the decay $H^\pm \rightarrow \bar{b}bW^*$ as a function of $m_{H^\pm}$ for $m_A = 100$ (solid lines) and 200 GeV (dotted lines) and three different values of $\tan \beta$.](image)

The situation described here corresponds to a particular direction of parameter space. One could have similarly allowed decays into $hW^+$ and $HW^+$. For instance, a search strategy based on tagging three $b$-quarks for each produced $t$-quark at the Tevatron (one $b$-jet coming form the $t \rightarrow bH^+$ decay and two $b$-jets coming from $H^+ \rightarrow W^+ + h, H, A$ with the Higgs bosons decaying into $b\bar{b}$ pairs) would then sum over all these decays. The corresponding theoretical branching ratio, however, becomes a function of $m_A$, $m_h$, $m_H$ and $\alpha$, in addition to $m_{H^\pm}$ and $\tan \beta$. Searches at LEP II and the Tevatron aimed at
constraining 2HDMs of Type II in the low tan β regime and/or 2HDMs of Type I will have to be modified accordingly. Constraints in the region of very large tan β for Type II couplings, when only the mode τ+ντ survives, remain unchanged.

In supersymmetric models, and in particular in the minimal version (MSSM), since m_A cannot be much smaller than m_{H^±} and the angle α is not an independent parameter, a non-trivial role is played only by the mode H^+ → hW^+. However, the branching ratio is large only for small values of the parameter tan β, tan β ≲ 2, for which the h boson is constrained to be rather heavy form LEP data [2] [in fact, such a low tan β scenario is by now excluded]. For larger values of tan β, the H^+Wh coupling is suppressed [and the H^+τν coupling is enhanced], making the branching ratio for this decay mode rather small, not exceeding ∼ 5% over the LEP allowed region. [Note that the situation might be different in extensions of the MSSM, such as in the case of additional singlet fields, the NMSSM, where m_{H^±} and m_A are not as strongly related as in the MSSM and the present LEP constraints on m_h and m_A do not hold; in this case BR(H^+ → hW^+, AW^+) might be rather large.]

In general, however, decays into the lightest chargino χ_1^+ and neutralino χ_1^0 as well as decays into sleptons are still allowed by present experimental data, and they dominate when they occur. (The importance of the channel χ_1^+χ_1^0 for a constrained minimal supersymmetric model was already discussed in [19]; for decays of MSSM Higgs bosons into supersymmetric particles, see also Ref. [20].)

The latest lower bounds on χ_1^+ from LEP II, m_{χ_1^+} ≳ 103.6 GeV, rely on the assumption of very heavy sleptons and/or a relatively large mass splitting with the lightest neutralino [21]. For large values of the Higgs–higgsino mass parameter μ, the lighter chargino and neutralino states χ_1^+ and χ_1^0 are respectively wino- and bino-like, with masses ∼ M_2 and ∼ M_1. In this case, even assuming gaugino mass universality at the very high scale: M_1 = \frac{5}{3}\tan^2θ_W M_2 ∼ \frac{1}{2}M_2, the decay channel H^+ → χ_1^+χ_1^0 is possible for m_{H^±} > 165 GeV. It gives rise to jets or leptons and missing energy and to τ’s and missing energy. The branching ratio BR(H^+ → χ_1^+χ_1^0) is shown in Fig. 3 as a function of m_{H^+}, for tan β = 4, M_2 = 150 GeV and μ = 200 GeV (solid line). [Here, and in the example for tan β = 4 in the next discussion, we have set the sfermion masses at ∼ 1 TeV and the trilinear stop coupling A_t at \sqrt{6} TeV (the so-called maximal mixing scenario) to evade the experimental bound [2] on the h bound mass.] For these values of parameters, χ_1^+ and χ_1^0 have respectively masses of 107 and 60 GeV.

The LEP II limits on χ_1^+ and χ_1^0 become weaker if the assumption on very heavy slepton masses and/or gaugino mass universality is relaxed. In both cases, the channel χ_1^+χ_1^0 becomes kinematically allowed for lighter H^±’s. As an example, we show in Fig. 3 the branching ratio in a direction of supersymmetric parameter space with M_1 disentangled from M_2 (dotted line). While keeping all other parameters fixed to the previous values, M_1 is set to 25 GeV, which induces a mass for χ_1^0 of ∼ 19 GeV. The mode χ_1^+χ_1^0 opens now already at ∼ 125 GeV. Figure 3 clearly shows that, in the region of moderate tan β, if no other decay of H^+ into superpartners is possible, the mode χ_1^+χ_1^0 can be dominant if it is kinematically allowed. For m_{H^±} ∼ 170 GeV and tan β = 4, the contribution of the χ_1^+χ_1^0 mode to the H^±’s total decay width, indeed, is respectively 78% and 92%
Figure 3: Branching fractions for the decay $H^+ \rightarrow \chi_1^+ \chi_1^0$ as a function of $m_{H^+}$, for $\tan \beta = 4$, $M_2 = 150$ GeV, $\mu = 200$ GeV and two different values of $M_1$: $M_1 \sim \frac{1}{2} M_2$ and $M_1 = 25$ GeV. All other supersymmetric decay modes are kinematically forbidden.

for $M_1 \sim M_2/2$ and $M_1 = 25$ GeV. An increase of $\tan \beta$ reduces the branching ratio $BR(H^+ \rightarrow \chi_1^+ \chi_1^0)$, while a smaller value of $\tan \beta$, if allowed, would make this decay mode even more dominant, in particular in the case of non–unified gaugino masses.

The existing lower bounds on the charged slepton masses from LEP II, are respectively 95, 88, and 76 GeV for $\tilde{e}$, $\tilde{\mu}$, $\tilde{\tau}$ when the mass difference with the lightest neutralino is rather large ($\Delta M \gtrsim 15$ GeV) and the sleptons are assumed to decay exclusively into $\ell^\pm \chi_1^0$ final states [22]. These bounds, in particular in the case of $\tilde{\tau}$, can be much weaker if they are nearly degenerate with the LSP neutralino. For sneutrinos, an absolute bound $\gtrsim 45$ GeV comes from the measurement of the invisible $Z$ boson decay width. Hence, the decay $H^+ \rightarrow \tilde{\tau}^+ \tilde{\nu}_\tau$ is therefore kinematically allowed and produces a final $\tau^+ +$ missing energy, but with a softer $\tau$ than that coming from the direct decay $H^+ \rightarrow \tau^+ \nu_\tau$. We show in Fig. 4 the relative branching ratio for two choices of input parameters:

a) $\tan \beta = 4$, $M_2 \sim 2 M_1 = 120$ GeV, $\mu = -500$ GeV, $m_{\tilde{\ell}_L} = m_{\tilde{\ell}_R} = m_{\tilde{\tau}} = 90$ GeV and $A_\tau = 0$ (small or moderate mixing scenario). This leads to a slepton spectrum: $m_{\tilde{\ell}} \sim 66$ GeV, $m_{\tilde{\mu}} \sim m_{\tilde{\tau}} \sim 100$ GeV and the two $\tilde{\tau}$ masses $\sim 20$ GeV below and above this value (the lightest chargino and neutralino masses are $m_{\chi_1^+} \sim 123$ GeV and $m_{\chi_1^0} \sim 60$ GeV).

b) $\tan \beta = 25$, $M_2 \sim 2 M_1 = \mu = 150$ GeV, $m_{\tilde{\ell}} = 100$ GeV and $A_\tau = -800$ GeV (the large $A_\tau$ value is chosen to maximize the $H^\pm \tilde{\tau} \tilde{\nu}_{\tau}$ coupling as will discussed later). This leads to the following spectrum: $m_{\tilde{\ell}} \sim 76$ GeV, $m_{\tilde{\mu}} \sim m_{\tilde{\tau}} \sim 110$ GeV and the $\tilde{\tau}$ mass $m_{\tilde{\tau}_1} \sim 63$ GeV almost degenerate with the lightest neutralino mass $m_{\chi_1^0} \sim 61$ GeV (therefore the decay $\tilde{\tau}_1 \rightarrow \chi_1^0 \tau$ gives very soft $\tau$ leptons, which will be overwhelmed by the $\gamma\gamma$ background and the LEP II lower limit on $m_{\tilde{\tau}_1}$ does not hold in this case).

Below the threshold for scenario a) with $\tan \beta = 4$, the dominant decays are $\tau^+ \nu_\tau$ and $hW^*$, while $AW^+$ and $c\bar{s}$ are below the percent level. Above the threshold, the branching ratio for the decay $H^+ \rightarrow \tilde{\tau}^+ \tilde{\nu}_\tau$ can become rather sizeable, possibly reaching the level
of $\sim 30\%$. For large enough $H^\pm$ masses, the channels $H^\pm \to \tilde{\mu}^\pm \tilde{\nu}_\mu$ and $\tilde{\epsilon}^\pm \tilde{\nu}_\epsilon$, open up, leading to an increase of $BR(H^\pm \to \tilde{\ell}\tilde{\nu})$ up to $\sim 80\%$. In scenario b) with $\tan\beta = 25$ and a large $A_\tau$ value, $\tau^+\nu_\tau$ decays are by far dominant below the threshold. When the decay $H^\pm \to \tilde{\tau}^\pm \tilde{\nu}_\tau$ opens up, the branching ratio quickly reaches the level of $\sim 75\%$.

The prominence of $\tau^+\nu_\tau$ decays observed above threshold is explained by the $H^\pm$ coupling to sleptons. For small stau mixing and small $\tan\beta$ values, the Lagrangian term $H^+\tilde{\nu}_L^*L, -(g/\sqrt{2})M_W\sin 2\beta$, is very large with respect to the Yukawa coupling $-(g/\sqrt{2})(m_\tau/M_W)\tan\beta$. Owing to the $\sin 2\beta$ dependence, this term quickly dies off for increasing $\tan\beta$. In this case, however, there exists other directions of parameter space where this decay mode still has a branching ratio $\sim 100\%$. For instance, when $A_\tau$ and $\tan\beta$ are large, since the coupling of the Lagrangian term $H^+\tilde{\nu}_L^*\tilde{\tau}_R: -(g/\sqrt{2})(m_\tau/M_W)(\mu + A_\tau\tan\beta)$ becomes very strong, the decay rate is enhanced as shown in Fig. 4 (note that for $A_\tau \sim \mu\tan\beta$, the left–right mixing in the slepton mass matrix tends to vanish).

![Figure 4: Branching fractions for the decay $H^+ \to \tilde{\ell}^\pm \tilde{\nu}_\ell$ as a function of $m_{H^+}$, for the two different sets of supersymmetric parameters a) and b) given in the text.](image)

Summarizing, at very large $\tan\beta$, a possible excess of $\tau$'s softer than those predicted by a 2HDM of Type II may indicate the presence of a heavier $H^\pm$ decaying into $\tilde{\tau}^+\tilde{\nu}_\tau$. Searches in the region of $\tan\beta \gtrsim 1$ should already consider multi-$b$ signals coming from $hW^+, bbW^+$ as well as $\tau$-signals with a wide momentum distribution coming from $\chi_1^+\chi_1^0$, $\tilde{\tau}^+\tilde{\nu}_\tau$, and $\tau^+\nu_\tau$ and jets/leptons + missing energy signals from $\chi_1^+\chi_1^0$.

It is needless to say that all these modes will play an important role in future searches and will not be blind to the intermediate range of $\tan\beta$. This would be particularly the case at the Tevatron Run II where the $H^\pm$ bosons, if light enough, can be produced copiously in top quark decays (other production channels would have much smaller rates).

While it would be always possible to detect them in a “disappearance” search (i.e. by looking at one top quark decaying into the standard mode, $t \to W^+b$, which should have a relatively large branching ratio, and ignoring the decay products of the other)
the direct search for 2HDMs $H^\pm$ bosons decaying into $Wb\bar{b}$ final states would be in principle relatively easy with high enough luminosity, since the performances of the CDF and D0 detectors for $b$–quark tagging are expected to be rather good. In the case of SUSY models, where the $H^\pm$ should be tagged though the leptonic decays of charginos or $\tau$ sleptons, the detection might be more challenging because of the softness of these particles. A detailed Monte–Carlo analysis, which is beyond the scope of this letter, will be needed to assess the potential of the Tevatron to search for the $H^\pm$ bosons in these new decay channels.

Note Added: After the first submission of this paper, a search for 2HDM charged Higgs bosons decaying into $AW^*$ finals states has been performed by the OPAL collaboration [24]; constraints in the $(m_A, m_{H^\pm})$ plane for various tan $\beta$ values have been set. In addition, the decay mode $H^\pm \rightarrow b\bar{b}W^\pm$ has been taken into account in simulations of $H^\pm$ searches at the upgraded Tevatron [23] and at the LHC [25]. Some of the decays modes discussed here have been also revisited in theoretical papers in the context of 2HDM [26] and the MSSM [27].

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