Charged and Pseudoscalar Higgs production at a Muon Collider

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Abstract

We consider single charged Higgs ($H^{\pm}$) and pseudoscalar Higgs ($A^0$) production in association with a gauge boson at $\mu^+\mu^-$ colliders. We find that the tree–level t–channel and s–channel contributions to $\mu^+\mu^- \rightarrow H^{\pm}W^{\mp}, A^0Z$ are enhanced for large values of $\tan \beta$, allowing sizeable cross–sections whose analogies at $e^+e^-$ colliders would be very small. These processes provide attractive new ways of producing such particles at $\mu^+\mu^-$ colliders and are superior to the conventional methods in regions of parameter space.
1 Introduction

Charged Higgs bosons ($H^\pm$) are predicted in many favourable extensions of the Standard Model (SM), in particular the Minimal Supersymmetric Standard Model (MSSM). Their phenomenology [1] has received much attention both at $e^+e^-$ colliders [2] and at hadron colliders [3, 4]. It is well known that $e^+e^-$ colliders offer a much cleaner environment in which to look than hadron colliders, although recently progress has been made in the possibilities of detecting $H^\pm$ for $M_{H^\pm} \geq m_t$ at hadron colliders [3]. At $e^+e^-$ colliders production proceeds via the mechanism $e^+e^- \to \gamma^*, Z^* \to H^+H^-$, with higher order corrections evaluated in [4], and detection is possible for $M_{H^\pm}$ up to approximately $\sqrt{s}/2$. The combined null–searches from all four LEP collaborations derive the lower limit $M_{H^\pm} \geq 77.3$ GeV (95% c.l) [7].

In recent years an increasing amount of work has been dedicated to the physics possibilities of $\mu^+\mu^-$ colliders [3, 4]. Such colliders offer novel ways of producing Higgs bosons, such as an $s$–channel resonance in the case of neutral scalars. The existing studies do not highlight any difference between the charged Higgs phenomenology at a $\mu^+\mu^-$ collider and $e^+e^-$ collider, and state that the main production mechanism would be via $\mu^+\mu^- \to \gamma^*, Z^* \to H^+H^-$. The rate for this process is identical at both colliders. In the MSSM $H^\pm$ becomes roughly degenerate in mass with $H^0$ and $A^0$ for masses greater than 200 GeV. It is this correlation among the masses of the Higgs bosons which disallows any large effects from a $s$–channel resonance (via $\mu^+\mu^- \to H^0, h^0 \to H^+H^-$) in the pair production mode, and we explicitly confirm this. In order for the above to be maximised one would require $\sqrt{s} \approx M_{h,H} \geq 2M_{H^\pm}$, a condition which requires sizeable mass splittings among the Higgs bosons and is disallowed in the MSSM.

So far unconsidered is the process $\mu^+\mu^- \to H^\pm W^\mp$ via $s$–channel and $t$–channel diagrams. Naively, this may offer greater possibilities of a large rate since the Yukawa coupling only appears at one vertex in contrast to both vertices in the pair production case. In addition, it offers the possibility of searching for $M_{H^\pm}$ up to $\sqrt{s} - M_W$ in contrast to pair production which only probes up to $M_{H^\pm} \leq \sqrt{s}/2$. The rate for $b\bar{b} \to H^\pm W^\mp$ at hadron colliders was considered in Ref. [11] although is not expected to provide an observable signature above the background [11], at least at LHC energies. In contrast, $\mu^+\mu^- \to H^\pm W^\mp$ might give a clean signature, since backgrounds are considerably less.

In an analogous way we also consider $\mu^+\mu^- \to A^0Z$. The phenomenology of $A^0$ is made tricky at $e^+e^-$ colliders due to the absence of a tree–level vertex $ZZA^0$ and so the standard Higgsstrahlung mechanism $(e^+e^- \to A^0Z)$ only proceeds via loops [12]. Moreover, over a wide region of parameter space in the MSSM $A^0$ has a suppressed rate in the channel $\mu^+\mu^- \to A^0h^0$, while $\mu^+\mu^- \to A^0H^0$ only probes up to $M_A \approx \sqrt{s}/2$. Proposed search strategies at $\mu^+\mu^-$ collider include the scanning technique and Bremsstrahlung tail method. Since both may provide a challenge for machine and detector design we consider the prospects of searching for $A^0$ via $\mu^+\mu^- \to A^0Z$.

Our work is organized as follows. In Section 2 we perform the full tree–level calculation of $\mu^+\mu^- \to H^+H^-, \mu^+\mu^- \to H^\pm W^\mp$ and $\mu^+\mu^- \to A^0Z$. In Section 3 we present numerical values of the cross–sections and Section 4 contains our conclusions.
2 Calculation

We now consider in turn the various production mechanisms. Our calculations are valid in both the MSSM and a general Two–Higgs–Doublet–Model (2HDM), the difference being that the MSSM Higgs sector is parametrized by just two parameters at tree–level (usually taken as $M_A$ and $\tan \beta$), while the 2HDM contains 7 free parameters. Thus in a general 2HDM all four Higgs boson masses may be taken as independent, as well as the two mixing angles $\alpha$ and $\beta$, and the Higgs potential parameter $\lambda_5$ (in the notation of Ref. [13]). In addition, the Higgs trilinear couplings differ from those in the MSSM. In this paper we shall present numerical results for the MSSM. Let us summarise the couplings needed for our study:

**Fermion–Fermion–Higgs couplings**

$$h^0 \mu^+ \mu^- = -\frac{igm^2}{2M_W} \lambda_{h^0 \mu^+ \mu^-}, \quad H^0 \mu^+ \mu^- = -\frac{igm^2}{2M_W} \lambda_{H^0 \mu^+ \mu^-}$$

$$A^0 \mu^+ \mu^- = -\frac{igm^2}{2M_W} \gamma_5 \lambda_{A^0 \mu^+ \mu^-}, \quad H^- \mu^+ \nu_\mu = \frac{igm^2}{\sqrt{2}M_W} \lambda_{H^+ \mu^+ \nu_\mu} \frac{1 - \gamma_5}{2} \quad (1)$$

In the MSSM these couplings are given by:

$$\lambda_{h^0 \mu^+ \mu^-} = \frac{\sin \alpha}{\cos \beta}, \quad \lambda_{H^0 \mu^+ \mu^-} = \frac{\cos \alpha}{\cos \beta}$$

$$\lambda_{A^0 \mu^+ \mu^-} = \tan \beta, \quad \lambda_{H^- \mu^+ \nu_\mu} = \tan \beta \quad (2)$$

One can see from the above formula that the CP–odd $A^0$ and the charged Higgs bosons coupling to the $\mu^\pm$ can be enhanced for large $\tan \beta$.

The momenta of the incoming $\mu^+$ and $\mu^-$, outgoing gauge boson $V$ ($W^\pm$ or $Z$) and outgoing Higgs scalar $S$ ($H^\pm$ or $A^0$) are denoted by $p_{\mu^+}$, $p_{\mu^-}$, $p_V$ and $p_S$, respectively. Neglecting the muon mass $m_\mu$, the momenta in the centre of mass of the $\mu^+ \mu^-$ system are given by:

$$p_{\mu^-, \mu^+} = \frac{\sqrt{s}}{2} (1, 0, 0, \pm 1)$$

$$p_{V, A^0} = \frac{\sqrt{s}}{2} (1 \pm \frac{M_V^2 - M_S^2}{s}, \pm \frac{1}{s} \lambda_4^\pm (s, M_V^2, M_S^2) \sin \theta, 0, \pm \frac{1}{s} \lambda_4^\pm (s, M_V^2, M_S^2) \cos \theta),$$

Here $\lambda(x, y, z) = x^2 + y^2 + z^2 - 2xy - 2xz - 2yz$ is the two body phase space function and $\theta$ is the scattering angle between $\mu^+$ and $S$; $M_V$ is the mass of the gauge boson $V$ and $M_S$ is the mass of the Higgs scalar $S$. In the case of $H^+ H^-$ production replace $V$ by $S$. The Mandelstam variables are defined as follows:

$$s = (p_{\mu^-} + p_{\mu^+})^2 = (p_V + p_S)^2$$

$$t = (p_{\mu^-} - p_V)^2 = (p_{\mu^+} - p_S)^2 = \frac{1}{2} (M_V^2 + M_S^2) - \frac{s}{2} + \frac{1}{2} \lambda_4^\pm (s, M_V^2, M_S^2) \cos \theta$$

$$u = (p_{\mu^-} - p_S)^2 = (p_{\mu^+} - p_V)^2 = \frac{1}{2} (M_V^2 + M_S^2) - \frac{s}{2} - \frac{1}{2} \lambda_4^\pm (s, M_V^2, M_S^2) \cos \theta$$

$$s + t + u = M_V^2 + M_S^2$$

$$3$$
2.1 \( \mu^+\mu^- \to H^+H^- \)

This process proceeds via the conventional Drell–Yan mechanism \( \mu^+\mu^- \to \gamma^*, Z^* \to H^+H^- \), the analogy of \( e^+e^- \to \gamma^*, Z^* \to H^+H^- \). Since \( m_\mu \approx 200m_e \) one may consider the s–channel and t–channel diagrams (see Fig. 1), whose analogies at \( e^+e^- \) colliders would be suppressed by factors of \( m_e \). The s–channel diagrams would be maximised for \( \sqrt{s} = M_h \) or \( M_H \), although in the context of the MSSM this condition would not allow on–shell pair production of \( H^\pm \). This can seen from the fact that \( \sqrt{s} \geq 2M_{h^\pm} \) and \( \sqrt{s} \approx M_h \) or \( M_H \) cannot be simultaneously satisfied in the MSSM. In contrast, such s–channel diagrams were considered in Ref.\[14\] for squark production via the process \( \mu^+\mu^- \to \tilde{q}\tilde{q} \), and were shown to cause a doubling of the cross–section at resonance. The t–channel diagram in Fig. 1 suffers from Yukawa coupling suppression at two vertices. In the calculation we shall use the following notation:

\[
Y_V = -Y_A = \frac{m_\mu^2}{4s_W^2M_W^2}\lambda_{h^\mu^\mu}\mu^\mu
\]

\[
a_h = -\frac{g_{HH^++\mu}m_\mu\lambda_{h^\mu^\mu}}{2M_Ws_W}, \quad a_H = -\frac{g_{HH^+H^-\mu}m_\mu\lambda_{h^\mu^\mu}}{2M_Ws_W}
\]

\[
a_1 = -\frac{1}{s} - \frac{1}{2s_Wc_W^2} - \frac{Y_V}{t}
\]

\[
a_2 = \frac{1}{2s_W^2c_W^2} - \frac{Y_A}{t}
\]

\[
a_3 = \frac{a_h}{s - M_W^2 + iM_W\Gamma_h} + \frac{a_H}{s - M_H^2 + iM_H\Gamma_H} + \frac{m_\mu Y_V}{t}
\]

where \( g_V = -(1 - 4s_W^2)/2 \), \( g_A = -1/2 \) and \( g_H = -c_W^2 + s_W^2 \). The coupling \( g_{HH^+H^-} \) and \( g_{HH^+H^-} \) (normalised to electric charge e) are given by:

\[
g_{HH^+H^-} = -\frac{1}{s_W}(M_W\cos(\beta - \alpha) - \frac{M_Z}{2c_W}\cos 2\beta \cos(\beta + \alpha) + \frac{c_\alpha \cos^2 \beta}{2c_W M_Z \sin \beta})
\]

\[
g_{HH^+H^-} = -\frac{1}{s_W}(M_W\sin(\beta - \alpha) + \frac{M_Z}{2c_W}\cos 2\beta \sin(\beta + \alpha) + \frac{\sin \alpha \cos^2 \beta}{2c_W M_Z \sin \beta})
\]

Where

\[
\epsilon = \frac{3G_Fm_t^4}{\sqrt{2\pi^2}\sin^2 \beta} \log \left[ \frac{m_t m_\mu}{m_t^2} \right]
\]
The $\epsilon$ term corresponds to the leading log 1–loop corrections [15] to the trilinear couplings. We will include also these leading log corrections to the Higgs–masses and to the mixing angles.

The square amplitude is given by:

$$|M|^2 = e^4 \left\{ \left( |a_1|^2 + |a_2|^2 \right) \frac{s^2}{2} \beta_H^2 \sin^2 \theta - 2|a_2|^2 m_H^2 s \beta_H + 2|a_3|^2 s \right. + 4 \Re(a_1 a_3) m_\mu s \beta_H \cos \theta \right\}$$  \hspace{1cm} (5)

with $\beta_H^2 = 1 - 4 M_H^2 s / s$. The differential cross–section is given by:

$$\frac{d\sigma}{d\Omega} = \frac{\beta_H}{64 \pi^2 s} |M|^2$$ \hspace{1cm} (6)

\subsection{2.2 $\mu^+ \mu^- \rightarrow H^\pm W^\mp$}

Single $H^\pm$ production may proceed via an s–channel resonance mediated by $h^0, H^0$ or $A^0$, and by t–channel exchange of $\nu_\mu$ (see Fig. 2). All are negligible at an $e^+ e^-$ collider due to the smallness of $m_e$. The loop induced contributions to $e^+ e^- \rightarrow H^\pm W^\mp$ were considered in Ref.[16] and shown to reach a few fb at very low values of $\tan \beta$, a region disfavoured in the MSSM. Potential advantages of $\mu^+ \mu^- \rightarrow H^\pm W^\mp$ over standard pair production are the following:

- $\mu^+ \mu^- \rightarrow H^\pm W^\mp$ is sensitive to the $H^\pm \mu^\mp \nu_\mu$ Yukawa coupling, which is model dependent, and hence provides information on the underlying Higgs structure. For example, we shall see that a 2HDM with the Model I type structure would not register a signal in this channel. In contrast $\mu^+ \mu^- \rightarrow \gamma^*, Z^* \rightarrow H^+ H^-$ has a model independent rate.

- Single $H^\pm$ production is less phase space suppressed than $H^\pm$ pair production, and would also allow greater kinematical reach at a given collider (on–shell production up to $\sim \sqrt{s} - M_W$).

- The t–channel contribution may be sizeable and does not require $\sqrt{s} \approx M_{res}$ to be significant, where $M_{res}$ is the mass of a neutral Higgs s–channel resonance. This is in contrast to other novel production processes at $\mu^+ \mu^-$ colliders, which usually require the condition $\sqrt{s} \approx M_{res}$. 

Figure.2
The differential cross-section for $\sigma(\mu^+\mu^- \rightarrow H^+W^\mp)$ may be written as follows:

$$\frac{d\sigma}{d\Omega} = \frac{\lambda^2(s, M_{H^\pm}^2, M_W^2)}{64\pi^2 s^2} |\mathcal{M}|^2$$

(7)

The matrix element squared is given by:

$$|\mathcal{M}|^2 = \frac{sg^4 m_\mu^2}{32 M_W^4} \left[ (|a_V|^2 + |a_A|^2) \lambda(s, M_{H^\pm}^2, M_W^2) + 2a_t^2 (2M_W^2 p_T^2 + t^2) + 2a_t(M_{H^\pm}^2 M_W^2 - s p_T^2 - t^2) \Re(a_V - a_A) \right]$$

(8)

Where $p_T^2 = \lambda(s, M_{H^\pm}^2, M_W^2) \sin^2 \theta/4s$ and the couplings $a_V, a_A$ and $a_t$ are given by:

$$a_V = \left( \frac{\cos(\alpha - \beta)\lambda_{h\mu^+\mu^-}}{s - M_h^2 + i M_h \Gamma_h} + \frac{\sin(\alpha - \beta)\lambda_{H\mu^+\mu^-}}{s - M_H^2 + i M_H \Gamma_H} \right)$$

(9)

$$a_A = \frac{\lambda_{A\mu^+\mu^-}}{s - M_A^2 + i M_A \Gamma_A}$$

(10)

$$a_t = \frac{\lambda_{H^-\mu^+\nu_\mu}}{t}$$

(11)

The mixing angle dependence of the Higgs–Fermion–Fermion couplings is contained in $\lambda_{h\mu^+\mu^-}, \lambda_{H\mu^+\mu^-}, \lambda_{A\mu^+\mu^-}$ and $\lambda_{H^-\mu^+\nu_\mu}$. Our formula agrees with that for $b\bar{b} \rightarrow H^\pm W^\mp$ in Ref. [10], with the replacements $m_t \rightarrow m_{\nu_\mu}$ and $m_b \rightarrow m_\mu$. Due to CP–invariance the rate for $W^+H^-$ and $W^-H^+$ production is identical. The total cross section takes the following form:

$$\sigma(\mu^+\mu^- \rightarrow W^+H^-) = \frac{G_F m_\mu^2}{16\pi s^2} \left( (|a_V|^2 + |a_A|^2) \lambda(s, M_{H^\pm}^2, M_W^2) s \right)$$

$$+ 2 \tan \beta [\Re(a_A - a_V)(M_{H^\pm}^2 + M_W^2 - s) + (s - 4M_W^2) \tan \beta \lambda^2(s, M_{H^\pm}^2, M_W^2)$$

$$- 4M_W^2 \tan \beta [\Re(a_V - a_A) M_{H^\pm}^2 s + (M_{H^\pm}^2 + M_W^2 - s) \tan \beta F(s, M_{H^\pm}^2, M_W^2)]$$

with:

$$F(s, M_S^2, M_T^2) = \log \left[ \frac{M_S^2 + M_T^2 - s - \lambda^2(s, M_S^2, M_T^2)}{M_S^2 + M_T^2 - s + \lambda^2(s, M_S^2, M_T^2)} \right]$$

(12)

2.3 $\mu^+\mu^- \rightarrow A^0Z$

As depicted in Fig. 3, this process proceeds in a very similar way to that for $\mu^+\mu^- \rightarrow H^\pm W^\mp$, except there are two t–channel diagrams. The process $\mu^+\mu^- \rightarrow Z\phi^0$, where $\phi^0$ is the SM Higgs boson, has been considered in Ref. [17]. Our calculation differs since there is no s–channel $Z$ exchange for $\mu^+\mu^- \rightarrow A^0Z$ in the MSSM. Instead there are two s–channel Higgs exchange diagrams of similar magnitude to the t–channel diagram, giving rise to strong interference. In addition $\tan \beta$ plays an important role. In the SM the s–channel $Z$ exchange is the dominant diagram at the collider energy we consider ($\sqrt{s} = 500$ GeV), and so interference is minimal.
The mechanism $\mu^+ \mu^- \rightarrow A^0 Z$ would provide an alternative way of searching for $A^0$ whose detection is not guaranteed at the LHC or a $\sqrt{s} = 500$ GeV $e^+ e^-$ collider. At the latter this is because the conventional production mechanism $e^+ e^- \rightarrow Z^* \rightarrow A^0 H^0$ would be closed kinematically for $M_A \approx M_H \geq 250$ GeV, and $e^+ e^- \rightarrow Z^* \rightarrow A^0 h^0$ ($\sim \cos^2(\beta - \alpha)$) is strongly suppressed for $M_A \geq 200$ GeV. The proposed search at a $\mu^+ \mu^-$ collider for $M_A \geq \sqrt{s}/2$ is by doing a scan over $\sqrt{s}$ energies, in order to find a resonance at $\sqrt{s} = M_A$, or by running the collider at full $\sqrt{s}$ and looking for peaks in the $b \bar{b}$ mass distribution (Bremsstrahlung tail method). These methods are competitive and both may allow detection up to $M_A \approx \sqrt{s}$ as long as $\tan \beta \geq 4 - 6$. However, both may provide quite a demanding challenge for detector resolution and machine design (see Ref. [8]), and it is too early to say with certainty if they would be feasible methods in practice. With this in mind we consider the process $\mu^+ \mu^- \rightarrow A^0 Z$. With a sizeable rate for $\sigma(\mu^+ \mu^- \rightarrow AZ)$, $A^0$ could be discovered first in this channel, and then the beams could be adjusted to $\sqrt{s} = M_A$ for precision studies. In addition, $\mu^+ \mu^- \rightarrow A^0 Z$ probes greater masses of $M_A$ than $e^+ e^- \rightarrow Z^* \rightarrow A^0 H^0$, and becomes another option to first discover $A^0$ (if discovery has been elusive at the LHC or a $\sqrt{s} = 500$ GeV $e^+ e^-$ collider).

The matrix element squared may be written as:

$$|\mathcal{M}|^2 = \frac{8g^4 m^2}{32M_W^2} \left[ |a_V|^2 \lambda(s, M_A, M_Z^2) - 2a_{t1}g_A(M_A^2 M_Z^2 - sp_T^2 - t^2) \Re(a_V) ight. $$

$$- 2a_{t2}g_A(M_A^2 M_Z^2 - sp_T^2 - u^2) \Re(a_V)$$

$$+ (g_A^2 + g_V^2) \left\{ a_{t1}^2 (2M_A^2 p_T^2 + t^2) + a_{t2}^2 (2M_Z^2 p_T^2 + u^2) \right\}$$

$$- 2(g_A^2 - g_V^2) a_{t1} a_{t2} (2M_Z^2 p_T^2 + 2M_A^2 M_Z^2 - tu) \right]$$

(13)

with $a_V$ the same as in Section 2.2 and

$$a_{t1} = \frac{\lambda_{A\mu^+ \mu^-}}{t - m_{\mu}^2}, \quad a_{t2} = \frac{\lambda_{A\mu^+ \mu^-}}{u - m_{\mu}^2}$$

(14)

The differential cross-section follows from eq(7) with the changes $M_{H^\pm} \rightarrow M_A$ and $M_W \rightarrow M_Z$.  

Figure 3
The total cross-section is given by:
\[
\sigma(\mu^+\mu^- \to A^0Z) = \frac{G_F m_e^2}{32\pi s^2} \left\{ [4s(\bar{g}_e^2 - g_\nu^2)\tan^2 \beta + 2s|a_V|^2\lambda(s, M_A^2, M_Z^2)]\lambda^\dagger(s, M_A^2, M_Z^2) \\
+ [8s\Re(a_V)g_A(M_A^2 + M_Z^2 - s)\tan \beta - 8(\bar{g}_e^2 - g_\nu^2)M_Z^2 \tan^2 \beta] \lambda^\dagger(s, M_A^2, M_Z^2) \\
+ 4(\bar{g}_e^2 + g_\nu^2)(s - 4M_Z^2)\tan^2 \beta] \lambda^\dagger(s, M_A^2, M_Z^2) \\
+ \frac{F(s, M_A^2, M_Z^2)}{(M_A^2 + M_Z^2 - s)}[-8M_Z^2 \tan \beta(-2\Re(a_V)g_A M_A^2 (M_A^2 + M_Z^2 - s)s \\
+ (2g_e^2 - g_\nu^2)M_A^2 (M_Z^2 - s) + (\bar{g}_e^2 + g_\nu^2)(M_A^2 + M_Z^2 - s)^2 \tan \beta)] \right\}
\] (15)

3 Numerical results

We now present our numerical analysis in the context of the MSSM. We take $\sqrt{s} = 500$ GeV and assume integrated luminosities of the order 50 fb$^{-1}$.

In Fig. 4 we plot $\sigma(\mu^+\mu^- \to H^\pm W^\mp)$, defined as the sum of $H^+W^-$ and $H^-W^+$ production, as a function of $M_{H^\pm}$, varying $\tan \beta$ from 20 to 50. We also include the tree-level rate for $\sigma(e^+e^- \to H^+H^-)$ in order to show the advantage of a $\mu^+\mu^-$ collider over an $e^+e^-$ collider. One can see that the single production mode gains in importance with increasing $\tan \beta$, and offers detection possibilities for $M_{H^\pm}$ up to $\sqrt{s} - M_W$. This compares favourably with the reach at an $e^+e^-$ collider.

The slight dip and rise of the curves arises due to the $H^0$ and $A^0$ mediated $s$-channel contributions increasing in magnitude with $M_{H^\pm}$, which compensates for the phase space suppression until the kinematical limit is approached. This can be seen from the fact that since $M_{H^\pm} \approx M_H \approx M_A$, larger $M_{H^\pm}$ causes both $M_H$ and $M_A$ to be closer to $\sqrt{s}$ (i.e. the resonance condition).

It is clear from the graphs that for $\tan \beta \geq 20$ one has $\sigma(\mu^+\mu^- \to H^\pm W^\mp) \geq 5$ fb, which would give a sizeable number of singly produced $H^\pm$ for luminosities of 50 fb$^{-1}$. One would expect $H^\pm \to tb$ decays for the mass region of interest and so the main background would be from $t\bar{t}$ production. Such a background was shown to overwhelm the channel $pp \to H^\pm W^\mp$ at the LHC. However, at a $\sqrt{s} = 500$ GeV muon collider $\sigma(\mu^+\mu^- \to t\bar{t}) \sim 0.7$ pb in contrast to $\sim 800$ pb at the LHC. Hence we would expect much better prospects for detection at a muon collider although a full signal-background analysis is beyond the scope of this paper. Previous studies of backgrounds to $H^\pm W^\mp$ production at $e^+e^-$ colliders have been carried out in the context of Higgs triplet models, assuming $H^\pm \to W^\pm Z$ as the main decay channel. Such studies cannot be applied to the MSSM where $H^\pm \to tb$ decays would dominate.

We note that a 2HDM with the Model I type structure would not register an observable signal in this channel. This is due to the rate being proportional to $\cot^2 \beta$, and so unacceptably small values of $\tan \beta$ would be required in order to allow observable cross-sections.

The process $\mu^+\mu^- \to A^0Z$ suffers from smaller cross-sections and these are plotted as a function of $M_A$ in Fig. 5. Given that $\mu^+\mu^- \to A^0H^0$ probes $M_A$ up to $\approx \sqrt{s}/2$ the
region $M_A \geq 250$ GeV is of interest. We see that cross–sections $\geq 1$ fb are only attainable in this region for $\tan \beta \geq 30$ and so detection would be restricted to large values of $\tan \beta$. The smallness of the cross–sections is caused by large destructive interference between the $s$ and $t$ channels. Finally, we consider $\mu^+\mu^- \rightarrow H^+H^-$. We find very small deviations from the rate for $e^+e^- \rightarrow H^+H^-$, of the order a few percent for large values of $\tan \beta$. This can be traced to the fact that the $s$–channel Higgs exchange diagrams are far from resonance, and the $t$–channel diagrams are doubly Yukawa suppressed. Since the 1–loop corrections may be much larger than these deviations we do not plot a graph.

4 Conclusions

We have considered the processes $\mu^+\mu^- \rightarrow H^\pm W^\mp$ and $\mu^+\mu^- \rightarrow A^0Z$ of the MSSM in the context of a high–energy $\mu^+\mu^-$ collider ($\sqrt{s} = 500$ GeV). We showed that $\mu^+\mu^- \rightarrow H^\pm W^\mp$ production offers an attractive new way of searching for $H^\pm$ at such colliders. The cross–section grows with increasing $\tan \beta$ with values as large as $30$ fb being attainable for $\tan \beta \geq 50$. With an integrated luminosity of $50$ fb$^{-1}$ a significant number of $H^\pm$ could be produced singly up to $M_{H^\pm} \approx \sqrt{s} - M_W$. This compares favourably with the reach
Figure 5: $\sigma(\mu^+\mu^- \rightarrow A^0Z)$ as a function of $M_A$ for various values of $\tan \beta$.

at an $e^+e^-$ collider, which may only probe up to $M_{H^\pm} \approx \sqrt{s}/2$. The main background (assuming $H^\pm \rightarrow tb$ decays) would be from $t\bar{t}$ production, which has a cross-section of 700 fb, 3 orders of magnitude less than at the LHC. We conclude that the mechanism $\mu^+\mu^- \rightarrow H^\pm W^\mp$ represents a novel and attractive way of producing $H^\pm$ at a $\mu^+\mu^-$ collider, and in our opinion merits a detailed signal–background analysis.

Pseudoscalar Higgs production via $\mu^+\mu^- \rightarrow A^0Z$ offers smaller cross-sections, with values of 2 fb or more only possible for large ($\geq 40$) $\tan \beta$. Charged Higgs pair production has essentially the same rate as that at an $e^+e^-$ collider, with differences of the order of a few percent for large values of $\tan \beta$.

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