Decays $h^{\pm} \to W^{\pm} h^0 (a^0)$ within an extension of the minimal supersymmetric standard model with one complex Higgs triplet

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The vertices $h^{\pm} W^{\mp} h^0$ and $h^{\pm} W^{\mp} a^0$, involving the gauge bosons $W^{\mp}$, the lightest charged ($h^{\pm}$), the lightest CP-even neutral (h^0), and the lightest CP-odd neutral (a^0) Higgs bosons, arise within the context of many extensions of the standard model, and they can be used to probe the Higgs sector of such extensions via the decays $h^{\pm} \to W^{\pm} h^0 (a^0)$. We discuss the strength of these vertices for an extension of the minimal supersymmetric standard model with an additional complex Higgs triplet. By using this model, we find regions of the parameter space where the decays $h^{\pm} \to W^{\pm} h^0 (a^0)$ are not only kinematically allowed, but they also become important decay modes and, in some cases, the dominant ones, with $\mathrm{BR}(h^{\pm} \to W^{\pm} h^0) \approx \mathrm{BR}(h^{\pm} \to W^{\pm} a^0)$.

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I. INTRODUCTION

The Higgs spectrum of many well motivated extensions of the standard model (SM) often include charged Higgs bosons whose detection in future colliders would constitute clear evidence of a Higgs sector beyond the minimal SM [1,2]. In particular, the two-Higgs-doublet model (THDM) has been extensively studied as a prototype of a Higgs sector that includes two charged Higgs bosons ($H^{\pm}$) [2]; however, a definitive test of the mechanism of electroweak symmetry breaking will require further studies of the complete Higgs spectrum. In addition, probing the properties of charged Higgs bosons through their decays could help find out whether they are indeed associated with a weakly interacting theory, as in the case of the popular minimal supersymmetric extension of the SM (MSSM) [3], or with a strongly interacting scenario [4]. Furthermore, these tests should also allow one to probe the symmetries of the Higgs potential and to determine whether the charged Higgs bosons belong to a weak doublet or to some larger multiplet.

Decays of charged Higgs bosons have been studied in the literature, including the radiative modes $W^{\pm}\gamma, W^{\pm}Z^0$ [5], mostly within the context of the THDM or its MSSM incarnation and, more recently, for the effective Lagrangian extension of the THDM [6]. Charged Higgs boson production at hadron colliders was studied long ago [7] and, recently, more systematic calculations of production processes at the future Large Hadron Collider (LHC) have been presented [8]. Current bounds on the mass of the charged Higgs bosons can be obtained at Tevatron, by studying the top decay $t \to bH^+$, which already eliminates some regions of the parameter space [9], whereas LEP-2 bounds give approximately $m_{H^+} > 80$ GeV [10].

On the other hand, the vertex $H^{\pm} W^{\mp} h^0$ deserves special attention because it can give valuable information about the underlying structure of the gauge and scalar sectors. In the first place, the decay mode $H^{\pm} \to W^{\pm} h^0$ might be detected at the LHC as claimed in Ref. [11], within the context of the MSSM. Furthermore, the vertex $H^{\pm} W^{\mp} h^0$ can also induce the associated production of $H^{\pm} h^0$ at hadron colliders, through a virtual $W^{\pm *}$ in the s channel which would become a relevant production mechanism for heavy charged Higgs bosons. In this paper, we are interested in studying the strength of this important vertex for an extension of the MSSM with one additional complex Higgs triplet (OHT-MSSM) [12,13], via the decay $h_k^{\pm} \to W^{\pm} h^0$, with $h^{\pm}$ the lightest charged and h^0 the lightest neutral Higgs bosons of the model.

This article is organized as follows: in Sec. II, we present and discuss briefly the results for the branching ratio (BR) of the charged Higgs boson decay in the context of the MSSM; we include in our numerical calculations the leading order radiative corrections. In Sec. III, we discuss the strength of the vertex for an extended supersymmetric model that includes a complex Higgs triplet (OHT-MSSM). We perform a numerical analysis to search for values of the Higgs boson masses that allow for the decay $H^{\pm} \to W^{\pm} h^0$. Finally, we summarize our conclusions in Sec. IV.

II. THE VERTEX $H^{\pm} W^{\mp} h^0$ IN THE MSSM

The simplest model that predicts charged Higgs bosons is the MSSM, which includes two scalar doublets of equal hypercharge, namely, $\Phi_1 = (\phi_1^+, \phi_1^0)$ and $\Phi_2 = (\phi_2^+, \phi_2^0)$. Besides two charged Higgs bosons ($H^{\pm}$), the spectrum of the MSSM includes two neutral CP-even states (h^0, H^0, with $m_{h^0} < m_{H^0}$), as well as a neutral CP-odd state (A^0).

Diagonalization of the charged mass matrices gives the expression for the charged Higgs boson mass eigenstates: $H^{\pm} = \cos\beta\phi_1^{\pm} + \sin\beta\phi_2^{\pm}$, where $\tan\beta(= v_2/v_1)$ denotes the ratio of vacuum expectation values (v.e.v.'s) of each doublet.

A. The decay $H^{\pm}\to W^{\pm}h^0$ in the MSSM with radiative corrections

Whenever kinematically allowed, the vertex $H^{\pm}W^{\pm}h^0$ could induce the decay $H^{\pm} \to W^{\pm}h^0$. For the light SM-like Higgs boson, this decay is proportional to the factor $\cos^2(\beta - \alpha)$, which determines its strength.

Other relevant decays of the charged Higgs boson are the modes into fermion pairs, which include the decays $H^{+(-)} \to \bar{\tau}\nu_\tau, c\bar{b}(\tau\bar{\nu}_\tau, \bar{c}b)$, and possibly into $t\bar{b}(\bar{t}b)$. If the charged Higgs bosons are indeed associated with the Higgs mechanism, their couplings to fermions should come from the Yukawa sector and the corresponding decays should have a larger BR for the modes involving the heavier fermions. The latter could be tested in a simple way if a comparison of the modes $H^{+(-)} \to \bar{\tau}\nu_\tau(\tau\bar{\nu}_\tau)$ and $H^{+(-)} \to \bar{\mu}\nu_\mu(\mu\bar{\nu}_\mu)$ led to very different BR's.

The masses of the two CP-even neutral Higgs bosons (h^0, H^0) and the charged pair $(H^{\pm})$ are conveniently determined in terms of the mass of the CP-odd state (A^0) and $\tan\beta$. In the MSSM the quartic couplings are given in terms of the gauge couplings, which implies that the light neutral Higgs boson must satisfy the (tree-level) bound $m_{h^0} \leq \cos 2\beta m_Z$. However, this relation is modified by important corrections arising from top/stop loops, which result into a bound $m_{h^0} \lesssim 130$ GeV [14].

In the decoupling limit $(m_A \gg m_Z)$ the parameters of the potential lead to the relation $\cos^2(\beta - \alpha) \simeq m_Z^2/m_{A^0}^2$, which remains small for large values of m_{A^0}. One also obtains an approximately degenerate spectrum of heavy Higgs bosons, i.e., $m_{H^{\pm}} \simeq m_{H^0} \simeq m_{A^0}$, while the mixing angles satisfy the following relation: $\alpha \simeq \beta - \pi/2$. Therefore, in the context of the MSSM, only the decay mode $W^{\pm}h^0$ is allowed for most regions of the parameter space. We have performed a detailed parametric search for contour regions for the branching ratio of $H^{\pm} \to W^{\pm}h^0$ by using the program HDECAY [15]; our results are shown in Fig. 1.

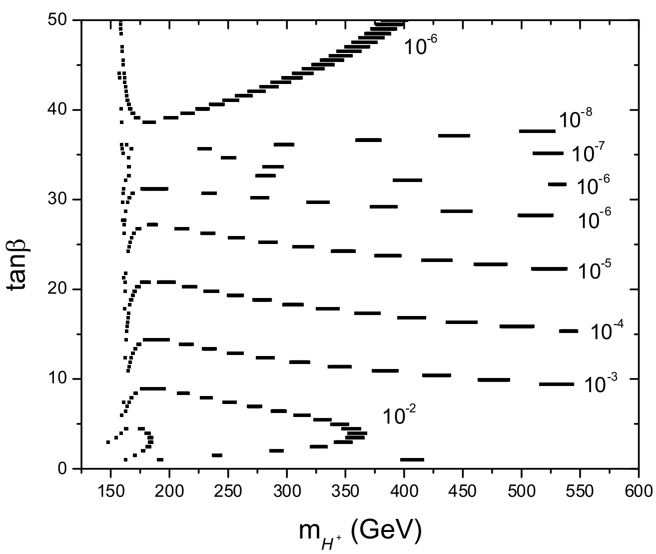

FIG. 1. BR $(H^{\pm} \to W^{\pm}h^0)$ in the MSSM, with radiative corrections to the Higgs mass as included in HDECAY, with $m_{\tilde{q}} = 500$ GeV, $\mu = 100$, and $A_0 = 1500$.

III. THE VERTEX $h^{\pm}W^{\mp}h^0$ IN A SUSY MODEL WITH AN ADDITIONAL COMPLEX HIGGS TRIPLET

The supersymmetric (SUSY) model with two doublets and a complex triplet (OHT-MSSM) [12,13] is one of the simplest extensions of the MSSM that allows one to study phenomenological consequences of an explicit breaking of the custodial symmetry SU(2) [13].

A. The Higgs sector of the model

The model includes two Higgs doublets and a (complex) Higgs triplet given by

$$\Phi_1 = \begin{pmatrix} \phi_1^0 \\ \phi_1^- \end{pmatrix}, \qquad \Phi_2 = \begin{pmatrix} \phi_2^+ \\ \phi_2^0 \end{pmatrix}, \\ \sum = \begin{pmatrix} \sqrt{\frac{1}{2}}\xi^0 & -\xi_2^+ \\ \xi_1^- & -\sqrt{\frac{1}{2}}\xi^0 \end{pmatrix}. \tag{1}$$

The Higgs triplet is described in terms of a 2×2 matrix representation; ξ^0 is the complex neutral field, and ξ_1^-, ξ_2^+ denote the charged scalars. The most general gauge invariant and renormalizable superpotential that can be written for the Higgs superfields $\Phi_{1,2}$ and Σ is given by

$$W = \lambda\Phi_1 \cdot \Sigma\Phi_2 + \mu_D\Phi_1 \cdot \Phi_2 + \mu_T \operatorname{Tr}(\Sigma^2), \tag{2}$$

where we have used the notation $\Phi_1 \cdot \Phi_2 \equiv \epsilon_{ab}\Phi_1^a\Phi_2^b$. The resulting scalar potential involving only the Higgs fields is thus written as

$$V = V_{\rm SB} + V_F + V_D,$$

where $V_{\rm SB}$ denotes the most general soft-supersymmetry-breaking potential [12]. In turn, the full scalar potential can be split into its neutral and charged parts, i.e., $V = V_{\rm charged} + V_{\rm neutral}$.

Besides the supersymmetry-breaking mass terms, m_i^2 $(i = 1, 2, 3)$, the potential depends on the parameters λ, μ_D, μ_T, A, B. For simplicity, we will assume that there is no CP violation in the Higgs sector, and thus, all the parameters and the v.e.v.'s are assumed to be real. The explicit expression of the Higgs potential is given in Ref. [12].

We can also combine the v.e.v.'s of the Higgs doublet as $v_D^2 \equiv v_1^2 + v_2^2$ and define $\tan\beta \equiv v_2/v_1$. Furthermore, the parameters v_D, v_T, m_W^2, and m_Z^2 are related as follows:

$$m_W^2 = \tfrac{1}{2}g^2(v_D^2 + 4v_T^2), \qquad m_Z^2 = \frac{\frac{1}{2}g^2 v_D^2}{\cos^2\theta_W}, \tag{3}$$

which implies that the ρ parameter is different from 1 at the tree level, namely,

$$\rho \equiv \frac{M_W^2}{M_Z^2\cos^2\theta_W} = 1 + 4R^2, \qquad R \equiv \frac{v_T}{v_D}. \tag{4}$$

The bound on R is obtained from the ρ parameter, which lies in the range 0.9993–1.0006, from the global fit reported in Refs. [14,16]. Thus, $R \le 0.012$ and $v_T \le 3$ GeV. We have taken into account this bound in our numerical analysis.

B. Mass spectrum

Diagonalization of the mass matrices (and the resulting mass eigenvalues) and mixing matrix will allow us to analyze the coupling $H_k^\pm W^\mp H_j^0$ ($k = 1, 2, 3$ and $j = 1, 2, 3$) and the coupling $H_k^\pm W^\mp A_j^0$ ($k = 1, 2, 3$ and $j = 1, 2$). The CP-even (odd) mass eigenstates are denoted by H_1^0, H_2^0, and H_3^0 (A_1^0 and A_2^0), ordered according to their masses, $m_{H_1^0} < m_{H_2^0} < m_{H_3^0}$ ($m_{A_1^0} < m_{A_2^0}$). The charged Higgs states are denoted by $H_k^\pm$ with $m_{H_1^\pm} < m_{H_2^\pm} < m_{H_3^\pm}$. We will denote the lightest charged scalar $H_1^\pm$, the lightest neutral scalar H_1^0, and the lightest neutral pseudoscalar A_1^0 as $h^\pm$, h^0, and a^0, respectively. Because of the large number of parameters appearing in our model, which include $\tan\beta$, R, λ, μ_D, μ_T, A, B_D, and B_T, it is convenient to consider only a few simple cases. In each, we will try to identify useful relations or trends for the behavior of the Higgs boson masses and couplings. In order to perform the numerical analysis leading to the allowed regions in the parameter space and the Higgs boson masses, we will make the following assumptions: (a) $\tan\beta$ is an independent variable; (b) R takes the representative value 0.01; (c) λ takes the value 0.5; and (d) the remaining parameters will cover the regions allowed by SUSY. Specifically, we will consider charged Higgs bosons masses in the range 100–300 GeV. Furthermore, we will restrict our numerical analysis to the following specific scenarios (which were introduced and discussed in Ref. [12]):

Scenario I $B_D = \mu_D = 0$, which represents the scenario when the spontaneous symmetry breaking (SSB) is dominated by the effects of the Higgs triplet, where we will consider the following cases: (A) $B_T = \mu_T = A$; (B) $B_T = \mu_T = -A$; (C) $B_T = -\mu_T = A$; (D) $-B_T = \mu_T = A$.

Scenario II $B_T = \mu_T = 0$. In this scenario the SSB is dominated by the effects of the Higgs doublets, where the following cases will be considered: (A) $B_D = \mu_D = A$; (B) $B_D = \mu_D = -A$; (C) $B_D = -\mu_D = A$; (D) $-B_D = \mu_D = A$.

Scenario III $|B_D| = |B_T| = |\mu_D| = |\mu_T| = |A|$. Both doublets and the triplet contribute to the SSB. Within this scenario several cases are considered: for instance, (A) $B_D = B_T = \mu_D = \mu_T = A$, as well as 15 other combinations with positive and negative signs.

For each point in the parameter space, within the above scenarios, we will determine the allowed regions by requiring the scalar squared mass eigenvalues to be positive and the Higgs potential lying in a global minimum. In these allowed regions, the masses of the physical Higgs bosons contained in the model are computed numerically.

C. The vertex $H_k^\pm W^\mp h^0$, $H_k^\pm W^\mp A^0$, and $H_k^\pm W^\mp Z^0$ ($k=1,2,3$)

We consider only the cases of the lightest neutral CP-even scalar, h^0, and the lightest neutral CP-odd scalar, a^0. We will use the expression for the vertex $H_k^\pm W^\mp h^0$ and $H_k^\pm W^\mp a^0$ for the OHT-MSSM reported in Refs. [12,13]. To present a complete study of the branching ratios of the charged Higgs bosons, we also discuss the vertex $H_k^\pm W^\mp Z^0$, which could dominate in some specific scenarios.

By using the expression for the rotation matrices of the charged and neutral Higgs sectors, U and V, we can write the coefficient of the vertex $H_k^\pm W^\mp h^0$ and $H_k^\pm W^\mp a^0$, namely, $\eta_k^{h^0}$ and $\eta_k^{a^0}$, respectively, as follows:

$$\eta_k^{h^0} = i\bigg(\frac{1}{\sqrt{2}}(V_{11}^S U_{2(k+1)} - V_{21}^S U_{1(k+1)}) + \frac{1}{4}V_{31}^S(U_{4(k+1)} - U_{3(k+1)})\bigg), \tag{5}$$

and

$$\eta_k^{a^0} = -i\bigg(\frac{1}{\sqrt{2}}(V_{11}^{PS} U_{2(k+1)} - V_{21}^S U_{1(k+1)}) + \frac{1}{4}V_{31}^{PS}(U_{4(k+1)} - U_{3(k+1)})\bigg), \tag{6}$$

where $H_k^\pm$ denote the charged Higgs bosons of the model, and h^0 (a^0) corresponds to the lightest neutral scalar (pseudoscalar) Higgs boson of the model. The U_{jk}'s denote the elements of the matrix, which relates the physical charged Higgs bosons (H_1^+, H_2^+, H_3^+) and the Goldstone boson G_0^+ (which gives mass to the W^+) with the fields: ϕ_2^+, ϕ_1^{-*}, ξ_2^+, and ξ_1^{-*}, as follows:

$$\begin{pmatrix} \phi_2^+ \\ \phi_1^{-*} \\ \xi_2^+ \\ \xi_1^{-*} \end{pmatrix} = \begin{pmatrix} U_{11} & U_{12} & U_{13} & U_{14} \\ U_{21} & U_{22} & U_{23} & U_{24} \\ U_{31} & U_{32} & U_{33} & U_{34} \\ U_{41} & U_{42} & U_{43} & U_{44} \end{pmatrix}\begin{pmatrix} G^+ \\ H_1^+ \\ H_2^+ \\ H_3^+ \end{pmatrix}. \tag{7}$$

The V_{ij}^S's and the V_{ij}^{PS}'s denote the elements of the rotation matrix for the CP-even and CP-odd neutral sectors, respectively. The matrices V^S and V^{PS} relate the physical scalars (H_1^0, H_2^0, H_3^0), the physical pseudoscalars (A_1^0, A_2^0), and Goldstone boson G^0 (which gives mass to the Z^0), with the real and imaginary parts of the fields ϕ_1^0, ϕ_2^0, ξ^0, in the following way:

$$\begin{pmatrix} \sqrt{\frac{1}{2}}\mathrm{Re}(\phi_1^0) \\ \sqrt{\frac{1}{2}}\mathrm{Re}(\phi_2^0) \\ \sqrt{\frac{1}{2}}\mathrm{Re}(\xi^0) \end{pmatrix} = \begin{pmatrix} V_{11}^S & V_{12}^S & V_{13}^S \\ V_{21}^S & V_{22}^S & V_{23}^S \\ V_{31}^S & V_{32}^S & V_{33}^S \end{pmatrix} \begin{pmatrix} H_1^0 \\ H_2^0 \\ H_3^0 \end{pmatrix} \qquad (8)$$

and

$$\begin{pmatrix} \sqrt{\frac{1}{2}}\mathrm{Im}(\phi_1^0) \\ \sqrt{\frac{1}{2}}\mathrm{Im}(\phi_2^0) \\ \sqrt{\frac{1}{2}}\mathrm{Im}(\xi^0) \end{pmatrix} = \begin{pmatrix} V_{11}^S & V_{12}^{PS} & V_{13}^{PS} \\ V_{21}^S & V_{22}^{PS} & V_{23}^{PS} \\ V_{31}^S & V_{32}^{PS} & V_{33}^{PS} \end{pmatrix} \begin{pmatrix} A_1^0 \\ G^0 \\ A_2^0 \end{pmatrix}. \qquad (9)$$

On the other hand, in this model the vertex $H_k^\pm W^\mp Z^0$ is also induced at tree level due to violation of the custodial symmetry. The expression for the vertex $H_k^\pm W^\mp Z^0$ is given by

$$H_k^\pm W_\mu^\mp Z_\nu^0: \ \pm i g^2 v_T (U_{3(k+1)} - U_{4(k+1)}) \cos\theta_W g_{\mu\nu}. \qquad (10)$$

One can see that only the triplet components contribute to this vertex, while the dependence on v_T gives a suppression effect.

D. Branching ratios for the principal two- and three-body decay modes of $h^\pm$

We now discuss the BR for the charged Higgs bosons, including the decay widths of the dominant modes of $h^\pm$, which turn out to be the following: (1) $h^\pm \to W^\pm Z^0$; (2) $h^\pm \to W^\pm h^0$; (3) $h^\pm \to W^\pm a^0$; (4) $h^{+(-)} \to t\bar{b}(\bar{t}b)$; (5) $h^{+(-)} \to \bar{\tau}\nu_\tau(\tau\bar{\nu})$. In order to discuss the BR for the charged Higgs bosons in the low mass region 100 GeV $< m_{h^\pm} <$ 200 GeV, it is necessary to include the dominant modes of the three-body decay of $h^\pm$, namely: (6) $h^\pm \to Z^0 W^{\pm *} \to Z^0 f\bar{f}'$; (7) $h^\pm \to h^0 W^{\pm *} \to h^0 f\bar{f}'$; (8) $h^\pm \to a^0 W^{\pm *} \to a^0 f\bar{f}'$ (it has been shown that this decay is a potentially strong tree-level process in the THDM-I [17,18]); (9) $h^+ \to t^* \bar{b} \to W^+ b\bar{b}$. The decay widths for each of the above modes are given as [19]:

(1) The decay $h^\pm \to W^\pm Z^0$:

$$\Gamma(h^\pm \to W^\pm Z^0) = g^2 v_T^2 |(U_{32} - U_{42})|^2 \cos^2\theta_W \lambda^{1/2}(1, \kappa_W, \kappa_Z) \left(\frac{(m_{h^\pm}^2 - m_W^2 - m_Z^2)^2 + 8 m_W^2 m_Z^2}{64\pi m_Z^2 m_W^2 m_{h^\pm}} \right). \qquad (11)$$

Here, $\kappa_W = m_W^2/m_{h^\pm}^2$ and $\kappa_Z = m_Z^2/m_{h^\pm}^2$, and $\lambda^{1/2}$ is the usual kinematic factor

$$\lambda^{1/2}(a, b, c) = \sqrt{(a - b - c)^2 - 4bc}. \qquad (12)$$

(2) The decay $h^\pm \to W^\pm h^0$:

$$\Gamma(h^\pm \to W^\pm h^0) = \frac{g^2 \lambda^{1/2}(m_{h^\pm}^2, m_W^2, m_{h^0}^2)}{64\pi m_{h^\pm}^3} |\eta_1^{h^0}|^2 \left[m_W^2 - 2(m_{h^\pm}^2 + m_{h^0}^2) + \frac{(m_{h^\pm}^2 - m_{h^0}^2)^2}{m_W^2} \right]. \qquad (13)$$

This decay is proportional to the factor $|\eta_1^{h^0}|^2$.

(3) The decay $h^\pm \to W^\pm a^0$:

$$\Gamma(h^\pm \to W^\pm a^0) = \frac{g^2 \lambda^{1/2}(m_{h^\pm}^2, m_W^2, m_{a^0}^2)}{64\pi m_{h^\pm}^3} |\eta_1^{a^0}|^2 \left[m_W^2 - 2(m_{h^\pm}^2 + m_{a^0}^2) + \frac{(m_{h^\pm}^2 - m_{a^0}^2)^2}{m_W^2} \right]. \qquad (14)$$

This decay is proportional to the factor $|\eta_1^{a^0}|^2$. In the MSSM the two-body decay of the charged Higgs boson into $W^\pm A^0$ is kinematically not allowed.

(4) The decay $h^{+(-)} \to t\bar{b}(\bar{t}b)$:

$$\Gamma(h^{+(-)} \to t\bar{b}(\bar{t}b)) = \frac{3g^2}{32\pi m_W^2 m_{h^\pm}^3} \lambda^{1/2}(m_{h^\pm}^2, m_t^2, m_b^2)[(m_{h^\pm}^2 - m_t^2 - m_b^2)(m_b^2 \tan^2\beta + m_t^2 \cot^2\beta) - 4 m_b^2 m_t^2]. \qquad (15)$$

(5) The decay $h^{+(-)} \to \bar{\tau}\nu_\tau(\tau\bar{\nu}_\tau)$:

$$\Gamma(h^{+(-)} \to \bar{\tau}\nu_\tau(\tau\bar{\nu}_\tau)) = \frac{g^2 m_\tau^2 \tan^2\beta}{32\pi m_W^2 m_{h^\pm}^3} (m_{h^\pm}^2 - m_\tau^2) \lambda^{1/2}(m_{h^\pm}^2, 0, m_\tau^2). \qquad (16)$$

(6) The decay $h^\pm \to Z^0 W^{\pm *} \to Z^0 f\bar{f}'$:

$$\Gamma(h^{\pm} \to Z^0 W^{\pm *} \to Z^0 f\bar{f}') = |F_Z|^2 \frac{3g^4 m_{h^{\pm}}}{512\pi^3} F(m_Z/m_{h^{\pm}}), \tag{17}$$

with $F_Z = (gv_T/m_W)(U_{32} - U_{42})\cos\theta_W$, and

$$F(x) = -|1-x^2|\left(\frac{47}{2}x^2 - \frac{13}{2} + \frac{1}{x^2}\right) - 3(1-6x^2+4x^4)|\ln x| + 3\frac{1-8x^2+20x^4}{\sqrt{4x^2-1}}\cos^{-1}\left(\frac{3x^2-1}{2x^3}\right).$$

We have simplified the expression for this width by taking the following approximation for the W propagator: $[(P-k)^2 - m_W^2]^{-1} \approx [m_{h^{\pm}}^2 - 2P\cdot k]^{-1}$, where P^{μ} and k^{μ} are the four momenta of the $H^{\pm}$ and Z^0 bosons, respectively.

(7) The decay $h^{\pm} \to h^0 W^{\pm *} \to h^0 f\bar{f}'$:

$$\Gamma(h^{\pm} \to h^0 W^{\pm *} \to h^0 f\bar{f}) = \frac{9g^4}{256\pi^3}|\eta_1^{h^0}|^2 m_{h^{\pm}} G_{h^0 W^{\pm}}, \tag{18}$$

where

$$G_{ij} = \frac{1}{4}\left\{2(-1+\kappa_j-\kappa_i)\sqrt{\lambda_{ij}}\left[\frac{\pi}{2} + \arctan\left(\frac{\kappa_j(1-\kappa_j+\kappa_i)-\lambda_{ij}}{(1-\kappa_i)\sqrt{\lambda_{ij}}}\right)\right] + (\lambda_{ij}-2\kappa_i)\log(\kappa_i) + \frac{1}{3}(1-\kappa_i)\left[5(1+\kappa_i) - 4\kappa_j - \frac{2}{\kappa_j}\lambda_{ij}\right]\right\} \tag{19}$$

and

$$\lambda_{ij} = -1 + 2\kappa_i + 2\kappa_j - (\kappa_i - \kappa_j)^2, \tag{20}$$

with $\kappa_i = m_i^2/m_{h^{\pm}}^2$.

(8) The decay $h^{\pm} \to a^0 W^{\pm *} \to a^0 f\bar{f}'$:

$$\Gamma(h^{\pm} \to h^0 W^{\pm *} \to h^0 f\bar{f}) = \frac{9g^4}{256\pi^3}|\eta_1^{a^0}|^2 m_{h^{\pm}} G_{a^0 W^{\pm}}. \tag{21}$$

The coefficient $G_{a^0W^{\pm}}$ has been defined in Eqs. (17) and (18).

(9) The decay $h^+ \to t^*\bar{b} \to W^+ b\bar{b}$:

$$\Gamma(h^+ \to t^*\bar{b} \to W^+ b\bar{b}) = \frac{1}{2}K_{H^{\pm}tb}\left\{\frac{\kappa_W^2}{\kappa_t^3}(4\kappa_W\kappa_t + 3\kappa_t - 4\kappa_W)\log\left(\frac{\kappa_W(\kappa_t-1)}{\kappa_t-\kappa_W}\right) + (3\kappa_t^2 - 4\kappa_t - 3\kappa_W^2 + 1)\log\left(\frac{\kappa_t-1}{\kappa_t-\kappa_W}\right) - \frac{5}{2} + \frac{1-\kappa_W}{\kappa_t^2}(3\kappa_t^2 - \kappa_t\kappa_W - 2\kappa_t\kappa_W^2 + 4\kappa_W^2) + \kappa_W\left(4 - \frac{3}{2}\kappa_W\right)\right\}, \tag{22}$$

with

$$K_{h^{\pm}tb} = \frac{3g^4 m_t^4}{1024\pi^3 m_W^4}\frac{1}{\tan^2\beta}m_{h^{\pm}}. \tag{23}$$

Finally, we evaluate numerically the BR for these nine principal modes, using the expressions given in Eqs. (9)–(21). In our computations, we will take $\sin^2\theta_W = 0.223$, $m_W = 80.4$ GeV, $m_b = 4.5$ GeV, $m_t = 174.3$ GeV [14]. We present in Figs. 2–10 our results in some specific scenarios, which can be summarized as follows:

In scenario I the numerical calculations are performed for $\lambda = 0.5$ and several values of $\tan\beta$. We consider case A within this scenario for the lightest charged Higgs boson $h^{\pm}$. We show our results for $\tan\beta = 15$, 30, and 50, in Figs. 2–4, respectively. In this scenario, $h^{\pm} \to W^{\pm}h^0$ and $h^{\pm} \to W^{\pm}a^0$ are the dominant decays for $\tan\beta = 15$. For $\tan\beta = 30$, the decays $h^{\pm} \to W^{\mp}h^0$ and $h^{\pm} \to W^{\pm}a^0$ are important modes, but $h^{\pm} \to W^{\pm}Z^0$ becomes the dominant decay. For $\tan\beta = 50$, $h^{\pm} \to W^{\pm}h^0$ and $h^{\pm} \to W^{\pm}a^0$ are still important decays (BR of the order of 10^{-1}), but we can see that $W^{\pm}Z^0$ is the dominant mode for $m_{h^{\pm}} < 195$ GeV and $t\bar{b}(\bar{t}b)$ is the dominant mode for $m_{h^{\pm}} > 200$ GeV.

In scenario II, where it mimicks the MSSM, we consider case D. Taking $\lambda = 0.5$, we plot our results for $\tan\beta = 5$, 15, and 30, in Figs. 5–7, respectively. In this scenario, $h^{\pm} \to W^{\pm}h^0$ and $h^{\pm} \to W^{\pm}a^0$ are important decays for 140 GeV $< m_{h^{\pm}} <$ 160 GeV, becoming the dominant decays for $\tan\beta \approx 5$.

Finally, for scenario III we consider case F ($B_D = B_T = -A$; $\mu_D = \mu_T = A$). Here both doublets and triplet contribute equally to the SSB. We calculate the BR's for the principal modes by taking $\lambda = 0.5$ and some values of $\tan\beta$. Our results are shown in Fig. 8 (for $\tan\beta = 5$), Fig. 9 (for $\tan\beta = 15$), and Fig. 10 (for $\tan\beta = 15$). The behavior of the BR of the different decay modes is similar to that observed in scenario IIA.

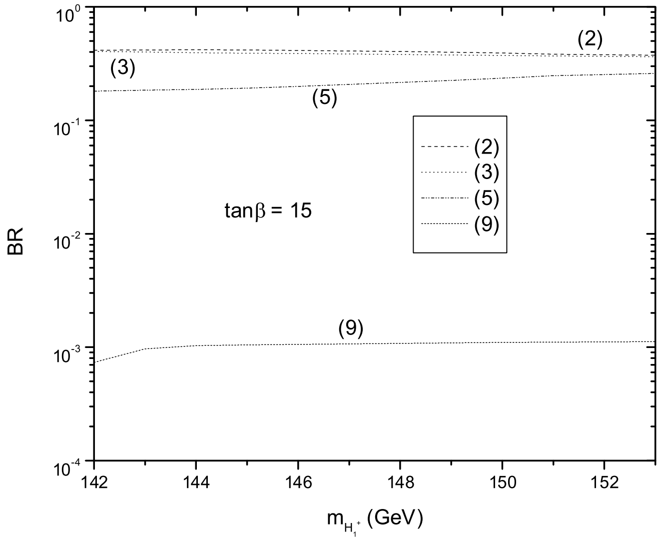

FIG. 2. Branching ratios of the charged Higgs bosons $h^{\pm}$ decaying into the principal modes for scenario I (case A), considering $\lambda = 0.5$. The various line drawings correspond to the different modes: (1) $h^{\pm} \rightarrow W^{\pm}Z^{0}$; (2) $h^{\pm} \rightarrow W^{\pm}h^{0}$; (3) $h^{\pm} \rightarrow W^{\pm}a^{0}$; (4) $h^{+(-)} \rightarrow t\bar{b}(\bar{t}b)$; (5) $h^{+(-)} \rightarrow \overline{\tau}\nu_{\tau}(\tau\overline{\nu})$; (6) $h^{\pm} \rightarrow Z^{0}W^{\pm *} \rightarrow Z^{0}f\overline{f}'$; (7) $h^{\pm} \rightarrow h^{0}W^{\pm *} \rightarrow h^{0}f\overline{f}'$; (8) $h^{\pm} \rightarrow a^{0}W^{\pm *} \rightarrow a^{0}f\overline{f}'$; (9) $h^{+} \rightarrow t^{*}\overline{b} \rightarrow W^{+}b\overline{b}$. These modes are shown for the lightest charged Higgs boson, for $\tan\beta = 15$.

To end this section, we want to point out the following. It is clear that for the OHT-MSSM there are regions in the parameter space that correspond to either the dominant (BR ≈ 1) or the moderate ($10^{-2} \lesssim$ BR $\lesssim 10^{-1}$) case. Therefore, the observation of the decays $h^{\pm} \rightarrow W^{\pm}h^{0}(a^{0})$, as the dominant modes, would back up the OHT-MSSM. On the other hand, the moderate case could arise of either the MSSM or the OHT-MSSM. The obser-

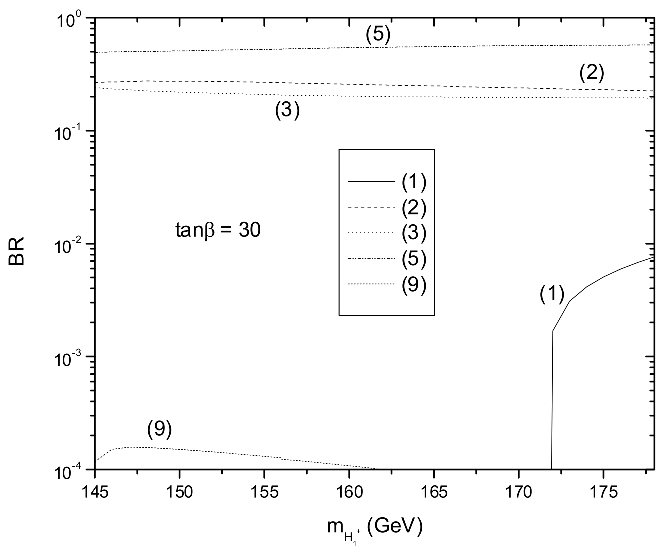

FIG. 3. Same as in Fig. 2, but for $\tan\beta = 30$.

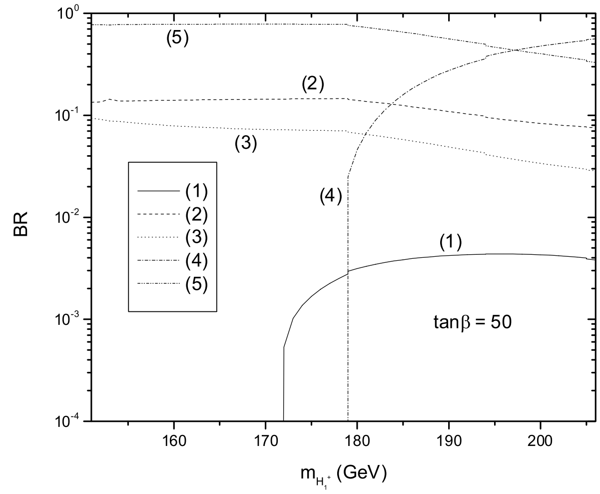

FIG. 4. Same as in Fig. 2, but for $\tan\beta = 50$.

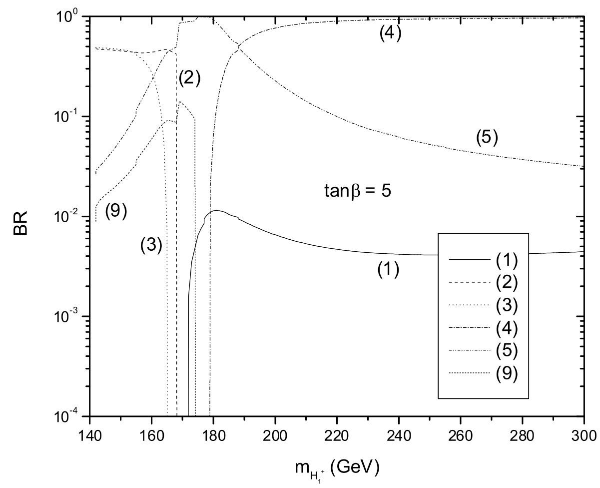

FIG. 5. Same as in Fig. 2, but for scenario II (case D), for $\tan\beta = 5$.

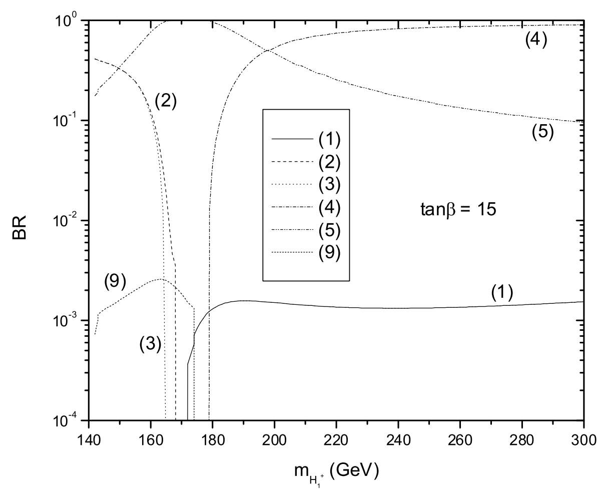

FIG. 6. Same as in Fig. 5, but for $\tan\beta = 15$.

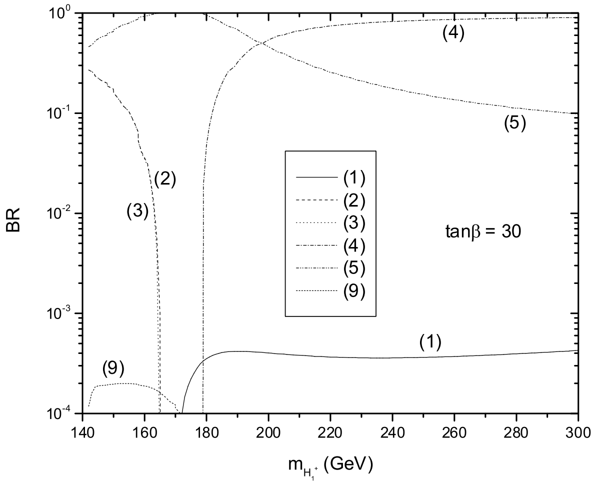

FIG. 7. Same as in Fig. 5, but for $\tan\beta = 30$.

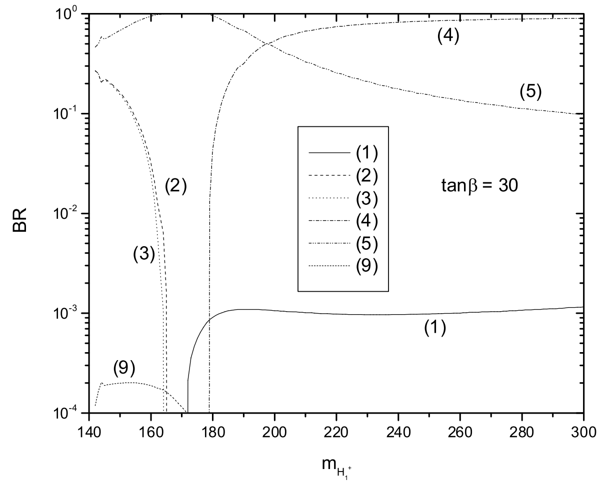

FIG. 10. Same as in Fig. 8, but for $\tan\beta = 30$.

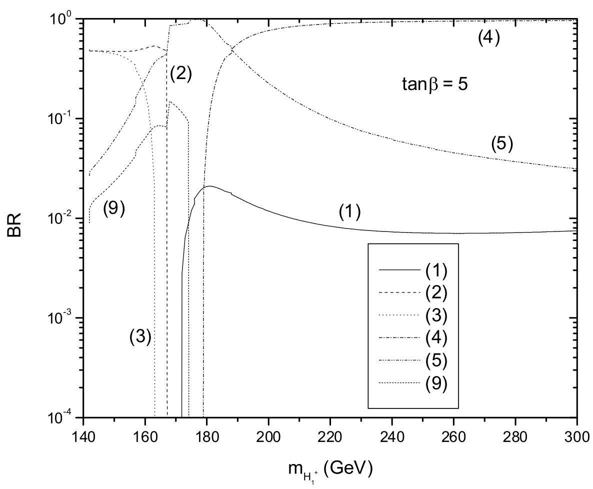

FIG. 8. Same as in Fig. 2, but for scenario III (case F), for $\tan\beta = 5$.

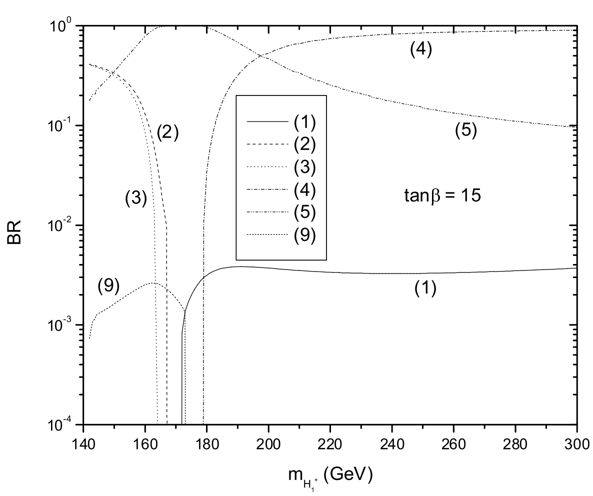

FIG. 9. Same as in Fig. 8, but for $\tan\beta = 15$.

vation of charged Higgs bosons in the region of the parameter space predicted by the MSSM would not discard the OHT-MSSM, while the detection of several charged Higgs bosons would correspond to a model with a more elaborate Higgs sector (such as Higgs triplets).

IV. CONCLUSIONS

We have studied the charged Higgs vertices $h^{\pm}W^{\mp}h^0(a^0)$, within the context of an extension of the minimal supersymmetric standard model with an additional complex Higgs triplet (OHT-MSSM) and then we have analyzed the decays $h^{\pm} \to W^{\pm}h^0(a^0)$ in the frame of this model. We found regions in the parameter space where the decays $h^{\pm} \to W^{\pm}h^0(a^0)$ are not only kinematically allowed, but they also become important decay modes and in some cases the dominant decay modes, with $\mathrm{BR}(h^{\pm} \to W^{\mp}a^0) \approx \mathrm{BR}(h^{\pm} \to W^{\mp}h^0)$. We conclude that for the OHT-MSSM there are regions in the parameter space that correspond to the case when the $W^{\pm}h^0(a^0)$ decay modes are dominant or gets a BR in the range 10^{-2}–10^{-1} (moderate case). The detection of the decay $h^{\pm} \to W^{\pm}h^0(a^0)$, as the dominant modes, would favor the SUSY triplet case. On the other hand, the moderate case could arise of either the MSSM or the OHT-MSSM. The detection of charged Higgs bosons in the region of the parameter space predicted by the MSSM would not discard the OHT-MSSM. Clearly, the observation of several charged Higgs bosons would correspond to a model with a more elaborate Higgs sector, such as the OHT-MSSM.

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