

# Charged Higgs boson in the $W^\pm$ Higgs channel at the Large Hadron Collider

Rikard Enberg<sup>a,\*</sup>, William Klemm<sup>a</sup>, Stefano Moretti<sup>b</sup>, Shoaib Munir<sup>a,c</sup>, Glenn Wouda<sup>a</sup>

<sup>a</sup> Department of Physics and Astronomy, Uppsala University, Box 516, SE-751 20 Uppsala, Sweden

<sup>b</sup> School of Physics & Astronomy, University of Southampton, Southampton SO17 1BJ, UK

<sup>c</sup> Asia Pacific Center for Theoretical Physics, San 31, Hyoja-dong, Nam-gu, Pohang 790-784, Republic of Korea

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## Abstract

In light of the recent discovery of a neutral Higgs boson,  $H_{\text{obs}}$ , with a mass near 125 GeV, we reassess the LHC discovery potential of a charged Higgs boson,  $H^\pm$ , in the  $W^\pm H_{\text{obs}}$  decay channel. This decay channel can be particularly important for a  $H^\pm$  heavier than the top quark, when it is produced through the  $pp \rightarrow t H^\pm$  process. The knowledge of the mass of  $H_{\text{obs}}$  provides an additional handle in the kinematic selection when reconstructing a Breit–Wigner resonance in the  $H_{\text{obs}} \rightarrow b\bar{b}$  decay channel. We consider some extensions of the Standard Model Higgs sector, with and without supersymmetry, and perform a dedicated signal-to-background analysis to test the scope of this channel for the LHC running at the design energy (14 TeV), for 300 fb<sup>-1</sup> (standard) and 3000 fb<sup>-1</sup> (high) integrated luminosities. We find that, while this channel does not show much promise for a supersymmetric  $H^\pm$  state, significant portions of the parameter spaces of several two-Higgs doublet models are testable.

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\* Corresponding author.

E-mail addresses: [Rikard.Enberg@physics.uu.se](mailto:Rikard.Enberg@physics.uu.se) (R. Enberg), [William.Klemm@physics.uu.se](mailto:William.Klemm@physics.uu.se) (W. Klemm), [S.Moretti@soton.ac.uk](mailto:S.Moretti@soton.ac.uk) (S. Moretti), [S.Munir@apctp.org](mailto:S.Munir@apctp.org) (S. Munir), [Glenn.Wouda@physics.uu.se](mailto:Glenn.Wouda@physics.uu.se) (G. Wouda).

## 1. Introduction

A charged Higgs boson,  $H^\pm$ , is predicted in many models of new physics, with and without Supersymmetry (SUSY). The observation of a  $H^\pm$  at the Large Hadron Collider (LHC) is thus expected to provide concrete evidence of physics beyond the Standard Model (SM). The strategies for such searches depend on the mass,  $m_{H^\pm}$ , of the charged Higgs boson. A  $H^\pm$  lighter than the top quark can be produced in  $t \rightarrow H^+ b$  and  $\bar{t} \rightarrow H^- \bar{b}$  decays, where the top quarks are produced in pairs in  $q\bar{q}$  annihilation and  $gg$  fusion (see [1] and the references therein). When  $m_{H^\pm} > m_t - m_b$ ,  $bg \rightarrow tH^-$  and  $gg \rightarrow tH^- \bar{b}$  are by far the dominant production processes.<sup>1</sup> As for the decays,  $H^\pm \rightarrow \tau\nu$ <sup>2</sup> is the dominant mode as long as  $m_{H^\pm} < m_t + m_b$ , beyond which  $H^\pm \rightarrow tb$  becomes the leading decay channel with branching ratio (BR) approaching unity.

The Minimal Supersymmetric Standard Model (MSSM) is an example of a scenario predicting charged Higgs states. In fact, it contains a total of five physical Higgs states. Among the neutral ones are included two CP-even states, with the lighter one denoted by  $h$  and the heavier by  $H$ , a CP-odd state,  $A$ , and there is also a charged pair  $H^\pm$ . The detection of an MSSM  $H^\pm$  lighter than the top quark is rather straightforward for a wide range of  $\tan\beta$  (where  $\tan\beta \equiv v_2/v_1$ , with  $v_1$  and  $v_2$  being the vacuum expectation values (VEVs) of the two Higgs doublet fields  $\Phi_1$  and  $\Phi_2$ ).  $H^\pm \rightarrow \tau\nu$  is the dominant decay mode of such a  $H^\pm$  for all  $\tan\beta$ . For  $m_{H^\pm} > m_t + m_b$ , the large reducible and irreducible backgrounds make the search for  $H^\pm$  in the  $tb$  decay mode notoriously difficult [10] (see [11,12] for experimental simulations). However, some studies [13, 14] concluded that the LHC discovery potential of a  $H^\pm$  state with mass  $\lesssim 600$  GeV is satisfactory in this decay channel, but only for very small,  $\lesssim 1.5$ , or very large,  $\gtrsim 30$ , values of  $\tan\beta$ . It has also been shown [15] that the  $H^\pm \rightarrow \tau\nu$  decay mode can be used at the LHC even for  $200 \text{ GeV} < m_{H^\pm} < 1 \text{ TeV}$  provided  $\tan\beta \gtrsim 3$ . In fact, if the distinctive  $\tau$ -polarisation [16] is used, the  $H^\pm \rightarrow \tau\nu$  channel can provide at least as good a heavy  $H^\pm$  signature as the  $H^\pm \rightarrow tb$  decay mode (for the large  $\tan\beta$  regime [17]).

At the LHC several searches have been carried out for  $H^\pm$ 's lighter as well as heavier than the top quark. The CMS collaboration has recently released exclusion limits [18] for a  $H^\pm$  lying in the 180 GeV–600 GeV mass range. That study assumes  $gg \rightarrow tH^- \bar{b}$  production and  $H^\pm \rightarrow tb$  and  $H^\pm \rightarrow \tau\nu$  decay modes and is based on  $19.7 \text{ fb}^{-1}$  of data collected at  $\sqrt{s} = 8 \text{ TeV}$ . An earlier analysis [19] based on the same dataset provided exclusion limits in the  $H^\pm \rightarrow \tau\nu$  decay channel for  $80 \text{ GeV} < m_{H^\pm} < 160 \text{ GeV}$ , assuming  $t\bar{t} \rightarrow H^\pm W^\pm b\bar{b}$  production, and for  $180 \text{ GeV} < m_{H^\pm} < 600 \text{ GeV}$ , using the inclusive  $pp \rightarrow tH^-(b)$  production mode. The same production and decay modes have also been analysed by the ATLAS collaboration [20] based on  $19.5 \text{ fb}^{-1}$  of data at  $\sqrt{s} = 8 \text{ TeV}$ , providing exclusion limits for  $80 \text{ GeV} < m_{H^\pm} < 160 \text{ GeV}$  and  $180 \text{ GeV} < m_{H^\pm} < 1 \text{ TeV}$ . In an earlier ATLAS study [21] based on  $4.7 \text{ fb}^{-1}$  of data at  $\sqrt{s} = 7 \text{ TeV}$ , the  $H^\pm \rightarrow cs$  decay channel has also been probed for  $H^\pm$  lying in the mass range 90 GeV–150 GeV.

<sup>1</sup> These are in fact one and the same process, describing the underlying dynamics in two different regimes, when combined with the parton distribution functions (pdfs). A combination of these two modes with a subtraction of the common terms is the preferred computational method, as described originally in [2,3] for neutral Higgs boson production and adapted later in [4,5] for charged Higgs boson production, with an implementation of the latter made available in [6,7]. (Also, see Refs. [8,9] for a discussion on the QCD accuracy at the next-to-leading order (NLO).) Further aspects in this context relevant to our analysis can be found in Section 5 below.

<sup>2</sup> We do not distinguish between fermions and anti-fermions when their identity is either unspecified or can be inferred from the context.

Note, however, that the two dominant decay channels mentioned above, i.e.,  $tb$  and  $\tau\nu$ , leave the  $1.5 \lesssim \tan\beta \lesssim 3$  window virtually unexplorable for a  $H^\pm$  heavier than the top quark in the MSSM. Importantly, it is for such small values of  $\tan\beta$  that the  $\text{BR}(H^\pm \rightarrow W^\pm h)$  becomes sizeable, reaching the percent level. The detectability of a Supersymmetric  $H^\pm$  in the  $W^\pm h$  decay channel was studied in [22], where it was noted that a  $H^\pm$  with mass around 200 GeV could be detectable at the LHC with  $\sqrt{s} = 14$  TeV and  $\mathcal{L} = 300 \text{ fb}^{-1}$ , for  $\tan\beta = 2\text{--}3$ . But there are two caveats. First, in these studies the mass of  $h$  was not fixed to the value eventually measured at the LHC. Second, such low values of  $\tan\beta$  may at first glance appear to be excluded by the LEP2 Higgs boson searches [23], particularly for low  $m_A \sim 100$  GeV. However, as discussed in [24], the LEP limit typically assumes a SUSY-breaking scale,  $M_{\text{SUSY}}$ , in the vicinity of 1 TeV, which should be relaxed owing to the fact that SUSY remains undiscovered, implying a significantly higher breaking scale. Now, a realistic SUSY model ought to contain a Higgs boson,  $H_{\text{obs}}$ , consistent with the one discovered at the LHC [25] and hence satisfying the ‘observational constraint,’  $122 \text{ GeV} \lesssim m_{H_{\text{obs}}} \lesssim 128 \text{ GeV}$ , which supersedes the LEP limit. The large allowed mass window is to take into account the theoretical uncertainties in the calculation of the  $H_{\text{obs}}$  mass in the model. All such aspects clearly need to be re-assessed in light of the latest experimental results.

Besides the above observational constraint on the mass of the Higgs boson, the LHC measurements of its signal strengths in various production and decay channels also strongly constrain the parameter space of the MSSM wherein a  $H^\pm$ , potentially visible via the  $W^\pm H_{\text{obs}}$  decay, can be obtained. In its singlet-extension, the Next-to-Minimal Supersymmetric Standard Model (NMSSM), the mass of the SM-like Higgs boson satisfying the mentioned mass constraint can be achieved in a more natural way, without requiring large radiative corrections from the stop sector. Such a Higgs boson, in fact, favours a lighter  $H^\pm$ , as we shall discuss in detail below. Moreover, in this model, which contains a total of 5 neutral Higgs states, the role of  $H_{\text{obs}}$  can be played by the any of the two lightest CP-even Higgs bosons,  $H_1$  or  $H_2$ , alternatively [26].

If one leaves aside SUSY, one of the simplest non-trivial extensions of the SM is represented by a 2-Higgs doublet model (2HDM), which contains two Higgs doublets with different Yukawa assignments (see [27] for a review). Notably, this structure (albeit limited to one specific Yukawa configuration) is necessary in the MSSM, implying that the Higgs spectrum in a CP-conserving 2HDM is the same as in the MSSM, containing three neutral Higgs bosons and a charged pair. However, the absence of SUSY relations amongst the Higgs boson masses allows much more freedom to alternatively identify the discovered SM-like Higgs state with either of the two CP-even Higgs bosons of a 2HDM. Depending on the way the Higgs doublets are assigned charges under a  $Z_2$  symmetry imposed in order to avoid large flavour-changing neutral currents (FCNCs), the 2HDMs are generally divided into four different types. In the ‘aligned’ 2HDM [28] (A2HDM), instead of the  $Z_2$  symmetry, a Yukawa-alignment is enforced in order to prevent large FCNCs.

From the point of view of  $H^\pm$  searches, results obtained in the MSSM can be easily translated to the case of a 2HDM Type II, as long as SUSY states are very heavy, i.e., decoupled [29]. This is somewhat more involved in the case of the other three ordinary Types and the A2HDM, although still possible (see [30] and [31], respectively). Some dedicated analyses of the 2HDMs to constrain them using the latest data from the LHC have also been performed recently [32]. The key phenomenological difference in the 2HDMs from the SUSY models in general, and the MSSM and NMSSM in particular, is that there are no light SUSY particles to provide cancellations (induced by the different spin statistics between SM and SUSY states) in low energy observables, chiefly from flavour dynamics. It is in fact the latter (e.g., limits on the  $Z \rightarrow b\bar{b}$  and

$b \rightarrow s\gamma$  decays) that generally produce severe constraints on the mass of  $H^\pm$  in the standard 2HDMs, pushing it to be larger than the top quark mass [33]. In the A2HDM, however, one can obtain  $m_{H^\pm} < m_t$  in a viable region of the parameter space [34].

In this article we analyse the possibility of establishing a  $H^\pm \rightarrow W^\pm H_{\text{obs}}$  signal in the next LHC run in all the models mentioned above, which are those where some relevance of such a decay has been established in the literature previously. We exploit the requirement on  $H_{\text{obs}}$  to have a mass around 125 GeV, so that the  $m_{H^\pm}$  range accessible via this signature starts at about 200 GeV and extends to nearly 500 GeV, as for heavier masses the  $tH^\pm$  production cross section becomes too low. We first discuss the consistency of the corresponding regions of the parameter spaces of these models with the current Higgs boson data from the LHC. We further assess the effects of imposing constraints from  $b$ -physics and, in the case of SUSY models, cold dark matter (DM) relic density measurements. We also carry out a model-independent detector-level analysis of the expected LHC sensitivity in the  $H^\pm \rightarrow W^\pm H_{\text{obs}}$  channel with  $\sqrt{s} = 14$  TeV. In doing so, we exploit the knowledge of the mass of  $H_{\text{obs}}$ , which will result in a substantial improvement in the efficiency of previously advocated [22] kinematical selections for the extraction of the signature of concern here, which we use for guidance. We then compare the sensitivities expected for various integrated luminosities at the LHC with the cross sections obtainable for this channel in each model considered in the presence of the aforementioned experimental constraints.<sup>3</sup> It will be the interplay between the improved selection and the reduced parameter space available following the Higgs boson discovery (with respect to the setups assumed in earlier analyses of the  $H^\pm$  decay mode considered here) that will determine the actual situation at present.

The article is organised as follows. In Section 2 we will discuss the production and decay mechanisms of the  $H^\pm$  considered in our analysis. In Section 3, we will discuss some salient features of the models analysed. In Section 4 we will provide some details of the scans of the parameter spaces of these models and of the experimental constraints imposed in our study. In Section 5 we will explain our signal-to-background analysis. In Section 6 we will present our results and in Section 7 our conclusions.

## 2. Production and decay of $H^\pm$

The dominant production process at the LHC for a  $H^\pm$  heavier than the top quark is its associated production with a single top, with the relevant subprocesses being  $bg \rightarrow tH^-$  and  $gg \rightarrow t\bar{b}H^-$  (plus charge conjugated channels). The division between these two subprocesses is not clear-cut. The  $gg$  amplitude can be seen as a tree-level contribution to the NLO amplitude that includes a virtual  $b$ -quark, with the  $bg$  process making the LO amplitude. In the  $gg$  process we may view the  $b$ -quarks (the virtual  $b$  and the emitted  $b$ ) as resulting from a splitting of the gluon and the corresponding amplitude contains the exact kinematics of this splitting. In the  $bg$  process the  $b$ -quark instead comes from the parton distribution of the proton. The  $b$ -quark is then a collinear parton arising from a splitting in the evolution of the pdfs. This contribution to the amplitude contains a collinear approximation of the kinematics and also a resummation of large logarithms in the factorisation scale that is not present in the  $gg$  amplitude.

When calculating the cross section for  $pp \rightarrow tH^\pm + X$  the  $bg$  and  $gg$  contributions to the amplitude cannot be added naively because that would result in double counting between the two contributions. There is a correct procedure to compute the total cross section [36], but it

<sup>3</sup> See [35] for a similar analysis for some Type II 2HDM benchmark points.

does not generalise to the differential cross section needed for Monte Carlo (MC) simulations. In Ref. [6] a method for event generation without double counting was introduced, and an add-on, called MATCHIG, to the event generator Pythia 6 [37] was constructed. In this framework events are generated both for the  $bg$  and  $gg$  processes and for the double-counting contribution. Events corresponding to the double counting have negative weights and should be subtracted from the positive weighted  $bg$  and  $gg$  processes. We have used MATCHIG in our simulations.<sup>4</sup>

The process  $pp \rightarrow tH^\pm + X$  has also been calculated at NLO and has been implemented [38] in the POWHEG BOX MC framework [39], which includes matching to parton showers. At NLO the  $bg$  and  $gg$  contributions are both part of the amplitude. It has also been implemented [40] in the MC@NLO framework [41]. In [38] it was shown that the MATCHIG program produces very similar kinematical distributions to the POWHEG implementation except at very large transverse momentum,  $p_T > 200$  GeV of the  $tH^\pm$  pair. The overall normalisation is, however, larger for the NLO calculations. The ratio between the total cross sections at NLO and LO depends on the model parameters via the mass spectrum, but for an example choice of 2HDMs it was found to be around a factor 2 for the Tevatron energies and a factor 1.4 for the LHC energies [38]. We do not consider this NLO enhancement of the signal in this paper for consistency, as we are only able to simulate the backgrounds at LO, but one should bear in mind that our quoted sensitivities may be somewhat stronger if NLO effects were systematically taken into account.

The spin/colour summed/averaged squared amplitude for the  $gb \rightarrow tH^-$  production process is given by [42]

$$|\overline{\mathcal{M}}|^2 = \frac{g_{qH^\pm}^2}{2m_W^2} \frac{g_s^2 g_2^2}{4N_c} |V_{tb}|^2 \frac{(u - m_{H^\pm}^2)^2}{s(m_t^2 - t)} \left[ 1 + 2 \frac{m_{H^\pm}^2 - m_t^2}{u - m_{H^\pm}^2} \left( 1 + \frac{m_t^2}{t - m_t^2} + \frac{m_{H^\pm}^2}{u - m_{H^\pm}^2} \right) \right], \quad (1)$$

where  $g_s$  and  $g_2$  are the  $SU(3)_C$  and  $SU(2)_L$  gauge couplings,  $N_c = 3$  is the number of colours and  $V_{tb}$  is the relevant CKM matrix element. See Refs. [14] and [43] for the  $gg \rightarrow tH^- \bar{b}$  amplitudes and graphs. The total cross section is proportional to the coupling  $g_{qH^\pm}^2$ , as noted in the equation above, which is the only model dependent factor for a given  $m_{H^\pm}$ . This factor depends on the masses,  $m_t$  and  $m_b$ , of the  $t$  and  $b$  quarks, respectively, as well as the parameter  $\tan\beta$ , and will be discussed in the next section for each model considered here. As shown in [6], the total cross section for a charged Higgs mass above  $m_t$  is actually well-approximated by the  $bg$  cross section. However, since the  $bg$  and the  $gg$  contributions lead to different kinematical distributions in the MC simulations, as noted above, we included both these contributions in our MC simulations.

Finally, as noted in the Introduction, this study aims to exploit the  $H^\pm \rightarrow W^\pm H_{\text{obs}}$  decay channel at the LHC. Of relevance for this particular process is the coupling of  $H^\pm$  to a generic neutral Higgs boson,  $H_i$ , and the  $W$  boson, given by

$$g_{H_i H^\pm W^\pm} = \frac{g_2}{2} (\cos\beta S_{i2} - \sin\beta S_{i1}), \quad (2)$$

where  $S_{i1}$  and  $S_{i2}$  are the elements of the mixing matrix that diagonalises the CP-even Higgs mass matrix in the model. It is clear that this coupling depends strongly on  $\tan\beta$ , both explicitly and through the elements  $S_{i1}$  and  $S_{i2}$ , (except in the A2HDM, as will be explained later) making the  $H^\pm \rightarrow W^\pm H_{\text{obs}}$  decay process highly sensitive to this parameter.

<sup>4</sup> The process  $bg \rightarrow tH^-$  already exists in the publicly available Pythia package.

### 3. The models

#### 3.1. Supersymmetric models

The Supersymmetric models considered here contain two Higgs doublets,  $\Phi_1$  and  $\Phi_2$ , which make the scalar components of the superfields  $\widehat{H}_d$  and  $\widehat{H}_u$ , respectively. The field  $\Phi_1$  is needed for generating the masses of the  $d$ -type quarks and leptons and  $\Phi_2$  those of the  $u$ -type quarks. The coupling of the charged Higgs boson to the quarks, defined in Eq. (1) as the factor  $g_{qH^\pm}^2$ , is given in these models as

$$g_{qH^\pm}^2 = m_b^2 \tan^2 \beta + m_t^2 \cot^2 \beta. \quad (3)$$

Thus the amplitude for the  $gb \rightarrow tH^-$  process is maximal for either small or large  $\tan \beta$ .

#### • MSSM

The MSSM Superpotential, from which the scalar potential is derived, is given as

$$W_{\text{MSSM}} = h_u \widehat{Q} \cdot \widehat{H}_u \widehat{U}_R^c + h_d \widehat{H}_d \cdot \widehat{Q} \widehat{D}_R^c + h_e \widehat{H}_d \cdot \widehat{L} \widehat{E}_R^c + \mu \widehat{H}_u \cdot \widehat{H}_d, \quad (4)$$

where  $\widehat{Q}$ ,  $\widehat{U}_R^c$ ,  $\widehat{D}_R^c$ ,  $\widehat{L}$  and  $\widehat{E}_R^c$  are the quark and lepton superfields and  $h_u$ ,  $h_d$  and  $h_e$  are the corresponding Yukawa couplings. In this model, the mass of  $H^\pm$  is given at LO as

$$m_{H^\pm}^2 = m_A^2 + m_W^2, \quad (5)$$

where  $m_W$  is the mass of the  $W$  boson. In order to allow the  $H^\pm \rightarrow W^\pm H_{\text{obs}}$  decay, one requires  $m_{H^\pm} > m_{H_{\text{obs}}} + m_W$ , which translates into the requirement  $m_A \gtrsim 190$  GeV. In the MSSM, under such a condition, the tree-level mass of the SM-like Higgs boson,  $H_{\text{SM}}$ , has an upper limit

$$m_{H_{\text{SM}}}^2 \leq m_Z^2 \cos^2 2\beta, \quad (6)$$

where  $m_Z$  is the mass of the  $Z$  boson. Therefore, if the  $H_{\text{SM}}$  is identified with the  $H_{\text{obs}}$  and hence required to have a mass close to 125 GeV in accordance with the LHC measurement, a large value of  $\tan \beta$  is necessary. Furthermore, the absence of any significant deviations of the signal strengths of the  $H_{\text{obs}}$  from the SM expectations so far [44] seems to be pushing the MSSM towards the so-called ‘decoupling regime’. This regime corresponds to  $m_A \gtrsim 150$  GeV for  $\tan \beta \gtrsim 10$  and yields SM-like couplings of the  $H_{\text{SM}}$ , in addition to a maximal tree-level mass, as noted above. The net effect of all these observations is that a  $H^\pm$  with mass greater than 200 GeV and a  $H_{\text{SM}}$  with the correct mass and SM-like couplings can be obtained simultaneously only for large  $\tan \beta$ . However, according to Eqs. (2) and (3),  $\tan \beta \sim 10$  not only diminishes the  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{SM}})$  but also the  $gb \rightarrow tH^-$  cross section.

The complete MSSM contains more than 120 free parameters in addition to those of the SM. In its phenomenological version, the pMSSM, one assumes the matrices for the sfermion masses and for the trilinear scalar couplings to be diagonal, which reduces the parameter space of the model considerably. Here, since we are mainly concerned with the Higgs sector of the model, we further impose the following mSUGRA-inspired (where mSUGRA stands for minimal supergravity) universality conditions:

$$\begin{aligned} m_0 &\equiv M_{Q_{1,2,3}} = M_{U_{1,2,3}} = M_{D_{1,2,3}} = M_{L_{1,2,3}} = M_{E_{1,2,3}}, \\ m_{1/2} &\equiv 2M_1 = M_2 = \frac{1}{3}M_3, \\ A_0 &\equiv A_t = A_b = A_\tau, \end{aligned} \quad (7)$$

where  $M_{Q_{1,2,3}}$ ,  $M_{U_{1,2,3}}$ ,  $M_{D_{1,2,3}}$ ,  $M_{L_{1,2,3}}$  and  $M_{E_{1,2,3}}$  are the soft masses of the sfermions,  $M_{1,2,3}$  those of the gauginos and  $A_{t,b,\tau}$  the soft trilinear couplings. This leaves us with a total of six free parameters, namely  $m_0$ ,  $m_{1/2}$ ,  $A_0$ ,  $m_A$ ,  $\tan\beta$  and the Higgs-higgsino mass parameter  $\mu$ .

### • NMSSM

The NMSSM [45–47] (see, e.g., [48,49] for reviews) contains a singlet Higgs field in addition to the two doublet fields of the MSSM. The scale-invariant Superpotential of the NMSSM is written as

$$W_{\text{NMSSM}} = \text{MSSM Yukawa terms} + \lambda \widehat{S} \widehat{H}_u \cdot \widehat{H}_d + \frac{\kappa}{3} \widehat{S}^3, \quad (8)$$

where  $\widehat{S}$  is the additional Higgs singlet Superfield and  $\lambda$  and  $\kappa$  are dimensionless Yukawa couplings. The introduction of the new singlet field results in a total of five neutral Higgs mass eigenstates and a  $H^\pm$  pair, after rotating away the Goldstone bosons. In the NMSSM, the MSSM upper limit on the tree-level mass of the SM-like Higgs boson, given in Eq. (6), gets modified as

$$m_{H_{\text{SM}}}^2 \leq m_Z^2 \cos^2 2\beta + \frac{\lambda^2 v^2 \sin^2 2\beta}{2} - \frac{\lambda^2 v^2}{2\kappa^2} \left[ \lambda - \sin 2\beta \left( \kappa + \frac{A_\lambda}{\sqrt{2}s} \right) \right]^2, \quad (9)$$

where  $v \equiv \sqrt{v_1^2 + v_2^2} = 246$  GeV,  $s$  is the VEV of the singlet field and  $A_\lambda$  is the soft SUSY-breaking parameter corresponding to the coupling  $\lambda$ . Clearly, for large values of  $\lambda$  and small  $\tan\beta$ , the second term in the above equation gives a significant positive contribution to the  $H_{\text{SM}}$  mass.

The mass expression for  $H^\pm$  in the NMSSM is given as

$$m_{H^\pm}^2 = m_A^2 + m_W^2 - \frac{v^2 \lambda^2}{2}, \quad (10)$$

where  $m_A^2$  is, in contrast with the MSSM, the diagonal entry  $[M_A^2]_{11}$  of the pseudoscalar mass matrix  $M_A^2$  of the model, given by

$$m_A^2 = [M_A^2]_{11} = \frac{\sqrt{2}\lambda s}{\sin 2\beta} \left( A_\lambda + \frac{\kappa s}{\sqrt{2}} \right). \quad (11)$$

Again, for a given value of  $\tan\beta$ , the negative third term in Eq. (10) results in a smaller  $m_{H^\pm}^2$  in the NMSSM compared to that in the MSSM, where it is given by the first two terms only. This negative contribution increases with the size of  $\lambda$ .

A crucial observation here is that a large  $\lambda$ , necessary to obtain sufficiently small  $m_{H^\pm}$ , has the dual advantage of enhancing also the tree-level mass of  $H_{\text{SM}}$ , as noted above. Such a scenario is therefore more natural than the one with a very MSSM-like  $H_{\text{SM}}$ , since a much smaller amount of fine-tuning is required to achieve the correct Higgs boson mass via radiative corrections. But large  $\lambda$  also implies a substantial singlet component in  $H_{\text{SM}}$ , which could result in significantly reducing its couplings to fermions and gauge bosons compared to those of the SM Higgs boson. However, recent studies [26] have shown that, for large  $\lambda$  and small  $\tan\beta$ , the  $H_{\text{SM}}$  of the model, which can correspond to either  $H_1$  or  $H_2$ , can still be consistent with the LHC Higgs boson data. The signal strength of  $H_{\text{SM}}$  in the  $\gamma\gamma$  decay channel in such a scenario can in fact be much larger than that of a SM-like Higgs boson, owing to a reduction in the  $\text{BR}(H_{\text{SM}} \rightarrow b\bar{b})$  compared to the true SM case. We point out here that, as in the MSSM, the  $H_{\text{SM}}$  in the NMSSM will also be identified with  $H_{\text{obs}}$ , since it is assumed to be the Higgs boson observed at the LHC.



The phenomenological version of the NMSSM that we study here contains three new parameters in addition to those of the pMSSM, mentioned earlier, with  $\mu$  replaced by  $\mu_{\text{eff}} (\equiv \lambda s)$  and  $m_A$  traded for  $A_\lambda$ . These include  $\lambda$ ,  $\kappa$  and  $A_\kappa$ , the latter being a dimensionful coupling originating in the SUSY-breaking part of the Higgs potential.

### 3.2. 2HDMs

A generic non-Supersymmetric 2HDM is defined by its scalar potential and its Yukawa couplings. The two Higgs doublets in such a model are written in terms of their VEVs and the physical Higgs states as

$$\Phi_1 = \frac{1}{\sqrt{2}} \begin{pmatrix} \sqrt{2}(G^+ \cos \beta - H^+ \sin \beta) \\ v_1 - h \sin \alpha + H \cos \alpha + i(G \cos \beta - A \sin \beta) \end{pmatrix}, \quad (12)$$

$$\Phi_2 = \frac{1}{\sqrt{2}} \begin{pmatrix} \sqrt{2}(G^+ \sin \beta + H^+ \cos \beta) \\ v_2 + h \cos \alpha + H \sin \alpha + i(G \sin \beta + A \cos \beta) \end{pmatrix}, \quad (13)$$

where  $\alpha$  is the mixing angle of the two CP-even Higgs bosons,  $\tan \beta$  has been defined earlier and  $G$  and  $G^+$  are the Goldstone bosons. The most general, CP-conserving potential for two Higgs doublets reads

$$\begin{aligned} \mathcal{V}_{2\text{HDM}} = & m_{11}^2 \Phi_1^\dagger \Phi_1 + m_{22}^2 \Phi_2^\dagger \Phi_2 - [m_{12}^2 \Phi_1^\dagger \Phi_2 + \text{h.c.}] \\ & + \frac{1}{2} \lambda_1 (\Phi_1^\dagger \Phi_1)^2 + \frac{1}{2} \lambda_2 (\Phi_2^\dagger \Phi_2)^2 + \lambda_3 (\Phi_1^\dagger \Phi_1) (\Phi_2^\dagger \Phi_2) + \lambda_4 (\Phi_1^\dagger \Phi_2) (\Phi_2^\dagger \Phi_1) \\ & + \left\{ \frac{1}{2} \lambda_5 (\Phi_1^\dagger \Phi_2)^2 + [\lambda_6 (\Phi_1^\dagger \Phi_1) + \lambda_7 (\Phi_2^\dagger \Phi_2)] \Phi_1^\dagger \Phi_2 + \text{h.c.} \right\}. \end{aligned} \quad (14)$$

Through the minimisation conditions of the Higgs potential above,  $m_{11}^2$  and  $m_{22}^2$  can be traded for the VEVs  $v_1$  and  $v_2$ , respectively. Furthermore, the tree-level mass relations allow the quartic coupling  $\lambda_{1-5}$  in Eq. (14) to be substituted by the four physical Higgs boson masses and the neutral mixing sector parameter  $\sin(\beta - \alpha)$ . Thus, in contrast with the SUSY models, in the 2HDMs the masses of the Higgs bosons are free input parameters, along with  $\lambda_6$ ,  $\lambda_7$ ,  $m_{12}^2$ ,  $\sin(\beta - \alpha)$  and  $\tan \beta$ .

In the 2HDMs, the Yukawa couplings of the fermions are also *a priori* free parameters. However, depending on how the two Higgs doublets couple to the fermions, FCNCs can be mediated by scalars at the tree level. The requirement of no large FCNCs thus puts very strong restrictions on the coupling matrices. There are two general approaches for avoiding large FCNCs. One way is to impose a  $Z_2$  symmetry so that each type of fermion only couples to one of the doublets (“natural flavour conservation”) [50,51]. The same symmetry then holds also in the scalar potential (forcing  $\lambda_6 = \lambda_7 = 0$ ), up to the soft breaking terms with parameter  $m_{12}^2$ , thus further reducing the number of free parameters.

As noted in the Introduction, there are four ways of assigning the  $Z_2$  charges, giving 2HDMs of Types I, II, X and Y. One defines as Type I the model where only the doublet  $\Phi_2$  couples to all fermions; Type II is the scenario similar to the MSSM, where  $\Phi_2$  couples to up-type quarks and  $\Phi_1$  couples to down-type quarks and leptons; in a Type X (or Type IV or ‘lepton-specific’) model  $\Phi_2$  couples to all quarks and  $\Phi_1$  couples to all leptons; and a Type Y (or Type III or ‘flipped’) model is built such that  $\Phi_2$  couples to up-type quarks and to leptons and  $\Phi_1$  couples to down-type quarks. The Type X and Type Y models have a similar phenomenology to Type I and II, respectively, especially in the context of this study. Specifically,  $g_{qH^\pm}^2$  is the same in the Type I and Type X models. Similarly, the Type Y model has a similar Yukawa structure, and



Table 1

The expressions for  $g_{qH^\pm}^2$  in the different 2HDMs considered in this paper.

	2HDM-I	2HDM-II	A2HDM
$g_{qH^\pm}^2$	$m_b^2 \cot^2 \beta + m_t^2 \cot^2 \beta$	$m_b^2 \tan^2 \beta + m_t^2 \cot^2 \beta$	$m_b^2 \tan^2 \beta^D + m_t^2 \tan^2 \beta^U$

consequently  $g_{qH^\pm}^2$ , as Type II, except for the leptons which couple to a different Higgs doublet in either of the two models. This, incidentally, implies that there is no  $\tan \beta$ -enhancement in the Type Y model to affect the  $\text{BR}(H^\pm \rightarrow \tau \nu)$ . We therefore consider only the Type I and Type II models, referred to as 2HDM-I and 2HDM-II, respectively, which are the most well-known ones.

Another way to achieve small FCNCs without imposing natural flavour conservation is to postulate that the Yukawa coupling matrices of the two Higgs doublets are proportional to each other, i.e., they are aligned. This approach has been adopted in the aforementioned A2HDM [28], where both scalar doublets ( $\Phi_1$  and  $\Phi_2$ ) couple to all types of fermions. In the  $Z_2$ -symmetric 2HDMs discussed above the Yukawa couplings are determined solely by the parameter  $\tan \beta$ , while the CP-conserving A2HDM instead has separate parameters for the up-type quarks, the down-type quarks and the leptons, usually denoted by  $\beta^U$ ,  $\beta^D$  and  $\beta^L$ . In the A2HDM there is no specific basis singled out by the fermionic sector due to the absence of the  $Z_2$  symmetry. For this study we choose the basis where only one doublet acquires a VEV, called the ‘Higgs basis’. In this basis the input parameters include  $\sin \alpha$  (where  $\alpha$  is the angle that diagonalises the CP-even Higgs-sector),  $\lambda_2$ ,  $\lambda_3$ ,  $\lambda_7$  and the above-mentioned alignment angles  $\beta^{U,D,L}$ , in addition to the physical Higgs boson masses.

The expressions for  $g_{qH^\pm}^2$  in Eq. (1) for the different 2HDMs (including the A2HDM) are given in Table 1. It should be noted that  $g_{qH^\pm}^2$  in the 2HDM-II is identical to the one in the SUSY models.

#### 4. Model scans and experimental constraints

We have performed scans of the parameter spaces of all the models considered here, requiring  $m_{H^\pm}$  to lie in the 200 GeV–500 GeV range. For each scenario except the MSSM, we carried out two separate scans for the cases with  $H_1$  and  $H_2$  alternatively playing the role of  $H_{\text{obs}}$ , i.e., having mass near 125 GeV and SM-like signal rates in the  $\gamma\gamma$  and  $ZZ$  decay channels. We point out here that in the MSSM it is not possible to obtain a  $H$  with a mass around 125 GeV while also requiring  $m_{H^\pm} \gtrsim 200$  GeV, as their masses lie very close to each other by theoretical construction. In the case of the SUSY models, since the masses of the scalar Higgs bosons are derived and not input parameters, we used the nested sampling package MultiNest-v2.18 [52] for efficiently scanning their parameter spaces.

The mass spectra and Higgs boson decay BRs for each scanned point of the MSSM, the NMSSM and the 2HDMs were computed using the public packages SUSY-HIT-v1.3 [53], NMSSMTools-v4.2.1 [54] and 2HDMC [55], respectively. For a point to be accepted in a given scan, it had to pass the condition  $122 \text{ GeV} \leq m_{H_{\text{obs}}} \leq 128 \text{ GeV}$  for the SUSY models and  $123 \text{ GeV} \leq m_{H_{\text{obs}}} \leq 127 \text{ GeV}$  in the 2HDMs. This is to take into account the experimental as well theoretical uncertainties (which are understandably larger in the presence of SUSY) in  $m_{H_{\text{obs}}}$  predicted in the two scenarios. As for the  $b$ -physics observables, the points for which their theoretically evaluated values did not lie in the following ranges were rejected during the scans for the NMSSM and the A2HDM.

- $2.63 \times 10^{-4} \leq \text{BR}(\bar{B} \rightarrow X_s \gamma) \leq 4.23 \times 10^{-4}$ ,
- $0.71 \times 10^{-4} < \text{BR}(B_u \rightarrow \tau \nu) < 2.57 \times 10^{-4}$ ,
- $1.3 \times 10^{-9} < \text{BR}(B_s \rightarrow \mu^+ \mu^-) < 4.5 \times 10^{-9}$ .

These 95% confidence level ranges are the ones suggested in the manual of the package SuperIso-v3.4 [56], which was used for the theoretical evaluation of these observables. Additionally, the scan points were also required to satisfy the constraint  $\Delta M_{B_d} = (0.507 \pm 0.004) \text{ ps}^{-1}$ , which is based on [57]. In the case of the  $Z_2$ -symmetric 2HDMs, their parameter spaces consistent with the  $b$ -physics constraints were adopted directly from [57], so that these constraints were not tested against during the scans. Moreover, for SUSY models the (lightest) neutralino DM relic density was calculated for every point using the package MicrOMEGAs-v2.4.5 [58]. Only points with  $\Omega_\chi h^2 < 0.131$ , assuming a +10% theoretical error on the central value of 0.119 measured by the PLANCK collaboration [59], were retained.

Finally, we used the public package HiggsBounds-v4.1.3 [60] to test the neutral Higgs bosons other than the  $H_{\text{obs}}$  in a given case for each model against the exclusion limits from the Large Electron–Positron (LEP) collider, the Tevatron and the LHC. This program also takes care of the exclusion constraints on  $H^\pm$  from the various LHC searches mentioned in the Introduction. Finally, the magnitude of a possible Higgs boson signal at the LHC is characterised by the signal strength modifier, defined as

$$\mu^X = \frac{\sigma(pp \rightarrow H_{\text{obs}} \rightarrow X)}{\sigma(pp \rightarrow h_{\text{SM}} \rightarrow X)}, \quad (15)$$

where  $X$  denotes the decay channel under consideration and  $h_{\text{SM}}$  denotes a 125 GeV SM Higgs boson. The theoretical counterparts of  $\mu^X$ , which we refer to as  $R^X$  here, were obtained from the program HiggsSignals-v1.20 [61] for  $X = \gamma\gamma, ZZ$ .<sup>5</sup> In our analysis below, while we will show all the good points from our scans, we will highlight the points for which  $R^{\gamma\gamma, ZZ}$  are consistent with the measured  $\mu^{\gamma\gamma, ZZ}$  at the LHC. The latest publicly available measurements read

$$\mu^{\gamma\gamma} = 1.13 \pm 0.24 \quad \text{and} \quad \mu^{ZZ} = 1.0 \pm 0.29 \quad (16)$$

at CMS [62] and

$$\mu^{\gamma\gamma} = 1.57^{+0.33}_{-0.28} \quad \text{and} \quad \mu^{ZZ} = 1.44^{+0.40}_{-0.35} \quad (17)$$

at ATLAS [63].<sup>6</sup>

## 5. Signal and background analysis

In addition to constraining the parameter spaces of the new physics models, knowledge of the mass of  $H_{\text{obs}}$  also provides an additional handle in identifying the  $H^\pm \rightarrow W^\pm H_{\text{obs}}$  decay. We focus here on the decay  $H_{\text{obs}} \rightarrow b\bar{b}$ , as it generally has a substantial BR and allows for a

<sup>5</sup> The  $\gamma\gamma$  and  $ZZ$  decay channels remain the only ones so far where a  $5\sigma$  excess has been established at the LHC.

<sup>6</sup> We note here that the ATLAS collaboration has recently made public [64] an updated measurement,  $\mu_{\gamma\gamma} = 1.17 \pm 0.27$ , which is now comparatively much closer to the SM prediction. However, no updates on  $\mu^{ZZ}$  for the same data set have been released. This implies that even if we use the newly released  $\mu^{\gamma\gamma}$  value, the older and larger value of  $\mu^{ZZ}$  in Eq. (17) will still rule out the corresponding model points, since  $R^{ZZ}$  is generally smaller than  $R^{\gamma\gamma}$ .

full reconstruction of  $H_{\text{obs}}$ .<sup>7</sup> In particular, we look for the production channel  $pp \rightarrow t(b)H^\pm \rightarrow W^\mp b(b)W^\pm H_{\text{obs}}$ , which, after semi-leptonic decays of the two  $W$  bosons and  $H_{\text{obs}} \rightarrow b\bar{b}$ , gives a final state of  $bbb(b)jj\ell\nu_\ell$ . The main background for this process is  $t\bar{t}$  production, and here we consider all processes  $pp \rightarrow t(b)W^\pm b\bar{b}$ , where the extra pair of  $b$ -quarks can come from the emission of a gluon, a Higgs boson, or a  $Z$ . In this section we describe our method for reconstructing the  $H^\pm$  signal and separating it from the background events to give an estimate of the sensitivities that could be achieved at the 14 TeV LHC.

We generate the hard process for the signal using the MATCHIG package [6] with Pythia 6.4.28 [37], thus including the  $bg$  and  $gg$  contributions and subtracting the correct double-counting term to get proper  $b$ -jet momentum distributions. Backgrounds were generated with MadGraph5 [65]. Parton showers and hadronisation for both signal and background were performed with Pythia 8 [66], followed by detector simulation with DELPHES 3 [67] using experimental parameters calibrated to the ATLAS experiment with modified  $b$ -tagging efficiencies.<sup>8</sup>

For reconstruction and background reduction, we roughly follow the procedures of previous analyses [22], with the addition of a top veto (described below) to further suppress the background.

1. Accept events with at least 3  $b$ -jets, at least 2 light jets, one lepton ( $e$  or  $\mu$ ), and missing energy. All objects must have transverse momentum  $p_T > 20$  GeV and rapidity  $|\eta| \leq 2.5$ , and must be separated from other objects by  $\Delta R > 0.4$ .
2. Find a hadronic  $W$  candidate from the light jets, taking the pair with the invariant mass  $m_{jj}$  closest to  $m_W$ . Reject the event if no pair satisfies  $|m_{jj} - m_W| \leq 30$  GeV.
3. Reconstruct a leptonically decaying  $W$  using the lepton and the missing energy, by assuming that the missing energy comes entirely from the single neutrino and imposing the invariant mass constraint  $m_{\ell\nu} = m_W$ . Because this is a quadratic constraint, there is a two-fold ambiguity in the solution for the longitudinal momentum of the neutrino. If the solutions are real, both are kept, and if they are complex, the real part is kept as a single solution.
4. Apply top veto for high mass searches (“veto first”).
5. Find a Higgs boson candidate from the  $b$ -jets, taking the pair with the invariant mass  $m_{bb}$  closest to  $m_{H_{\text{obs}}} \approx 125$  GeV. Reject the event if no pair satisfies  $|m_{bb} - m_{H_{\text{obs}}}| \leq 15$  GeV.
6. Apply top veto for low mass searches (“veto second”).
7. Reconstruct a top quark using the remaining  $b$ -tagged jet(s) and reconstructed  $W$ ’s, taking the combination which gives  $m_{bW}$  closest to  $m_t$ . If one of the leptonically-decaying  $W$  solutions is selected here, the other is discarded. Reject the event if no combination satisfies  $|m_{bW} - m_t| \leq 30$  GeV.
8. Reconstruct the charged Higgs candidate from the remaining  $W$  and the reconstructed  $H_{\text{obs}}$  to determine the discriminating variable  $m_{WH_{\text{obs}}}$ .

Because the largest background is by far  $t\bar{t}X$ , we wish to suppress it as much as possible by identifying events in which a top quark pair can be reconstructed. The majority of  $t\bar{t}X$  events

<sup>7</sup> This channel was also recently studied in [35], where it was noted that especially when uncertainties become dominated by systematics, the decay  $H_{\text{obs}} \rightarrow \tau^+\tau^-$  can become more relevant due to its smaller backgrounds, despite a smaller BR and additional unobservable neutrinos. In this study, we consider only statistical uncertainties.

<sup>8</sup> The  $b$ -tagging used is given by  $\epsilon_\eta \tanh(0.03 p_T - 0.4)$ , with the transverse momentum,  $p_T$ , in GeV,  $\epsilon_\eta = 0.7$  for central ( $|\eta| \leq 1.2$ ), and  $\epsilon_\eta = 0.6$  for forward ( $1.2 \leq |\eta| \leq 2.5$ ) jets. This choice is a conservative one in comparison with the ATLAS high-luminosity projections [68].

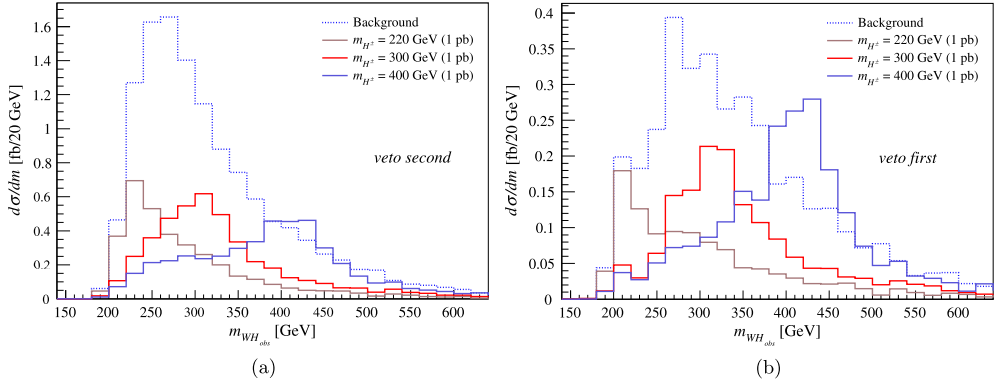


Fig. 1. Reconstructed  $m_W H_{\text{obs}}$  for signal and background with two different top vetos: (a) first identify an  $H_{\text{obs}} \rightarrow b\bar{b}$  candidate, then veto event if two top jets can be reconstructed with remaining objects (veto second); (b) using all final state objects, veto event if two top jets can be reconstructed (veto first). The signal is normalised to  $\sigma(pp \rightarrow tH^\pm) \times \text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}}) \times \text{BR}(H_{\text{obs}} \rightarrow b\bar{b}) = 1$  pb before selection and cuts.

which are able to pass our requirement of providing an SM-like Higgs candidate do so by combining a  $b$ -jet coming from a top decay with another  $b$ -tagged jet, so the background will be most reduced if a top veto is applied before the Higgs reconstruction,

**Veto first:** Using reconstructed  $W$ 's and all remaining jets, veto event if two top quarks can be reconstructed, both with  $|m_{Wj} - m_t| \leq 20$  GeV.

We also wish to avoid unnecessarily cutting signal events. When a charged Higgs boson with  $m_{H^\pm} \geq m_t$  undergoes the decay  $H^\pm \rightarrow W^\pm H_{\text{obs}} \rightarrow W^\pm b\bar{b}$ , it is kinematically possible for one of the  $b$ -jets from the  $H_{\text{obs}}$  decay to combine with the  $W$  to give an invariant mass close to the top mass. Indeed, this effect occurs in large regions of the available phase space for charged Higgs bosons with masses just above the threshold for  $W^\pm H_{\text{obs}}$  decays. In this case, we wish to identify the  $b\bar{b}$  pair from the  $H_{\text{obs}}$  decay before applying a top veto,

**Veto second:** After identifying two  $b$ -jets which reconstruct  $H_{\text{obs}}$ , using reconstructed  $W$ s and all remaining jets, veto event if two top quarks can be reconstructed, both with  $|m_{Wj} - m_t| \leq 20$  GeV.

Figs. 1(a) and 1(b) show the signal and background  $m_W H_{\text{obs}}$  distributions for  $m_{H^\pm} = 220, 300, 400$  GeV and the two types of top veto. The “veto first” scenario clearly reduces the background more effectively, but at the expense of a reduced signal. However, for larger  $m_{H^\pm}$ , the signal is less likely to fake an additional top, so there is less difference between the two vetoes in the higher mass signal distributions.

It is also clear from Fig. 1 that the  $H^\pm$  resonance can be reconstructed well enough to further separate it from the background. For each mass, we select a window in the reconstructed  $m_W H_{\text{obs}}$  range which maximises the statistical significance  $S/\sqrt{B}$  of the signal.<sup>9</sup> We additionally choose

<sup>9</sup> In events where a leptonic  $W$  with two real solutions is used in the reconstruction, the event is accepted if either solution gives an  $m_W H_{\text{obs}}$  within the window.

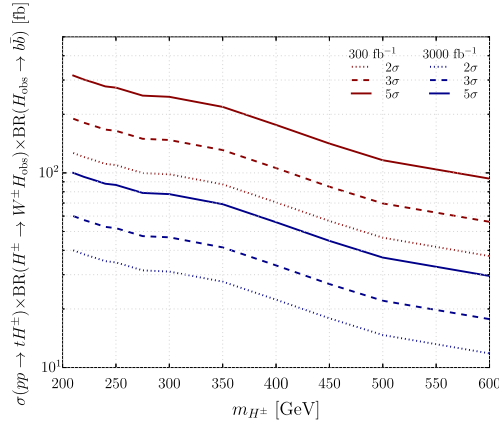


Fig. 2. Sensitivity of the LHC to the signal cross section for exclusion, evidence and discovery, based on statistical uncertainties. Contours are thus shown for  $S/\sqrt{B} = 2, 3, 5$  for an integrated luminosity of  $\mathcal{L} = 300 \text{ fb}^{-1}$  at the next LHC run and at the high luminosity LHC with  $\mathcal{L} = 3000 \text{ fb}^{-1}$ , both at  $\sqrt{s} = 14 \text{ TeV}$ .

the top veto which maximises  $S/\sqrt{B}$  for each mass, and find that “veto second” is most effective at lower masses,  $m_{H^\pm} \lesssim 350 \text{ GeV}$ , whereas “veto first” is preferable above this mass range.<sup>10</sup> In Fig. 2 we show how this signal and background translate into sensitivities at the 14 TeV LHC for different values of the product  $\sigma(pp \rightarrow tH^\pm) \times \text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}}) \times \text{BR}(H_{\text{obs}} \rightarrow b\bar{b})$ , which we henceforth refer to as the signal cross section. We see that we can probe  $\sigma \times \text{BR} \sim \mathcal{O}(100 \text{ fb})$  with an integrated luminosity of  $300 \text{ fb}^{-1}$ , but require higher luminosities to see  $\mathcal{O}(10 \text{ fb})$  signals. These sensitivities can be compared to the model-dependent cross sections and BRs in various scenarios, which we discuss in the following section.

## 6. Results and discussion

### 6.1. MSSM

In Fig. 3(a) we show the mass of  $h$  as a function of  $m_{H^\pm}$  in the MSSM, with the heat map corresponding to  $\tan\beta$ . The ranges of the MSSM input parameters scanned to obtain these points are shown in Table 2(a). One sees in the figure that for the selected  $m_{H^\pm}$  range,  $m_{H_{\text{SM}}}$  lying between 122 GeV–128 GeV can only be obtained for  $\tan\beta \gtrsim 6$ . As noted earlier, such intermediate values of  $\tan\beta$  bring down not only the  $pp \rightarrow tH^\pm$  cross section but also the  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$ . The product of these two quantities, only for points in the narrow strip corresponding to  $m_{H_{\text{SM}}} > 122 \text{ GeV}$  and consequently to highest allowed  $\tan\beta$  in Fig. 3(a), is shown in Fig. 3(b). This product hardly exceeds 4 fb, and that too only for points very close to the lower limit imposed on  $m_{H_{\text{SM}}}$ . The heat map in the figure shows the  $\text{BR}(H_{\text{obs}} \rightarrow b\bar{b})$ , which grows as the  $H_{\text{obs}}$  becomes more and more SM-like due to falling  $m_A$ , and hence  $m_{H^\pm}$ , given the intermediate value of  $\tan\beta$ .

<sup>10</sup> As already mentioned, here we consider only statistical uncertainties (and give the significance as  $S/\sqrt{B}$ ). A full experimental analysis with all errors included might prefer a different mass for the transition between vetoes.

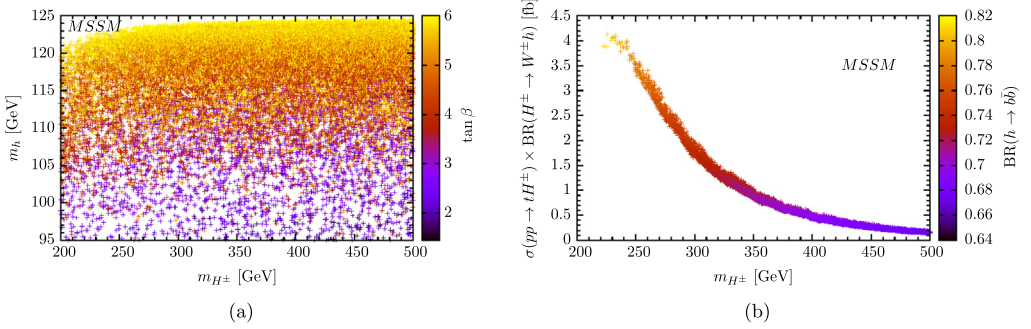


Fig. 3. (a)  $m_h$  as a function of  $m_{H^\pm}$  in the MSSM, with the heat map showing the parameter  $\tan \beta$ . (b)  $\sigma(pp \rightarrow t H^\pm) \times \text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  as a function of  $m_{H^\pm}$  in the MSSM, with the heat map showing the  $\text{BR}(H_{\text{obs}} \rightarrow b\bar{b})$ .

Table 2

Ranges of the input parameters scanned for (a) the MSSM and (b) the NMSSM.

MSSM parameter	Range	NMSSM parameter	Range
$m_0$ (GeV)	500–4000	$m_0$ (GeV)	500–3000
$m_{1/2}$ (GeV)	300–2000	$m_{1/2}$ (GeV)	300–2000
$A_0$ (GeV)	–7000–7000	$A_0$ (GeV)	–4000–4000
$\mu$ (GeV)	100–2000	$\tan \beta$	1–6
$m_A$ (GeV)	100–500	$\lambda$	0.45–0.7
$\tan \beta$	1–6	$\kappa$	0.2–0.5
		$\mu_{\text{eff}}$ (GeV)	100–200
		$A_\lambda$ (GeV)	0–500
		$A_\kappa$ (GeV)	–500–0

## 6.2. NMSSM

Our initial scans for the NMSSM covered very wide ranges of the nine input parameters mentioned in Section 3. These scans revealed only a small region of the NMSSM-specific parameters where  $m_{H_{\text{obs}}}$  and  $m_{H^\pm}$  both lied within the desired ranges. Two subsequent scans of this narrow region, for the cases with  $H_{\text{obs}} = H_1$  and with  $H_{\text{obs}} = H_2$  each, yielded a much larger density of interesting points. The corresponding parameter ranges are given in Table 2(b).

In Fig. 4(a) we show the  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  as a function of  $m_{H^\pm}$  for the points obtained in the scan requiring  $H_1$  to be the  $H_{\text{obs}}$ . In Fig. 4(b) the corresponding points for the case with  $H_{\text{obs}} = H_2$  are shown. The heat maps in the two figures show the distribution of the  $\sigma(pp \rightarrow t H^\pm)$ . We see in the figures that while the  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  in the  $H_1 = H_{\text{obs}}$  ( $H_2 = H_{\text{obs}}$ ) case can reach up to  $\sim 23\%$  ( $\sim 28\%$ ), its maximum reachable value drops slowly with decreasing  $m_{H^\pm}$  and, in fact, for  $m_{H^\pm} < 250$  GeV it falls below 5%. This behaviour of the  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  is thus in conflict with that of the  $\sigma(pp \rightarrow t H^\pm)$ , which clearly rises with decreasing  $m_{H^\pm}$  and is in fact maximal for points with the lowest  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  observed.

In Fig. 5(a) we show the signal cross section for the case with  $H_{\text{obs}} = H_1$ . The points in green are the ones fulfilling only the  $b$ -physics constraints and we note for these points that, as a result of the tension between the  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  and the  $\sigma(pp \rightarrow t H^\pm)$ , the total cross section barely exceeds 10 fb. The points in red and blue in the figure are the ones for which  $R^{\gamma\gamma/ZZ}$  are consistent with the CMS and ATLAS ranges of  $\mu^{\gamma\gamma/ZZ}$ , respectively. Evidently, imposing



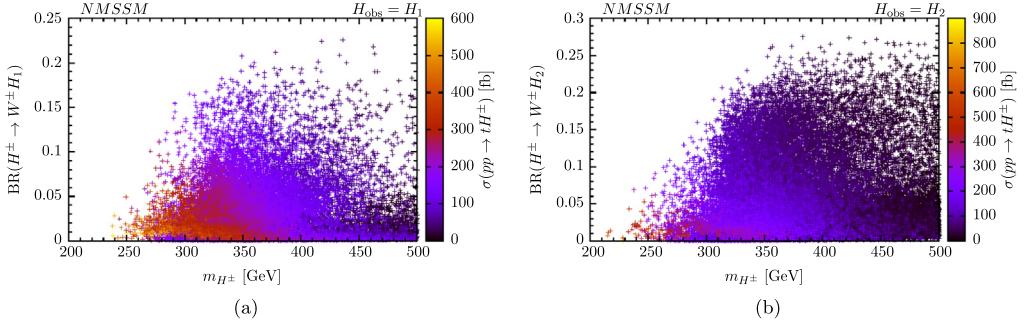


Fig. 4.  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  as a function of  $m_{H^\pm}$  in the NMSSM when (a)  $H_{\text{obs}} = H_1$  and (b)  $H_{\text{obs}} = H_2$ , with the heat map showing the  $\sigma(pp \rightarrow tH^\pm)$ .

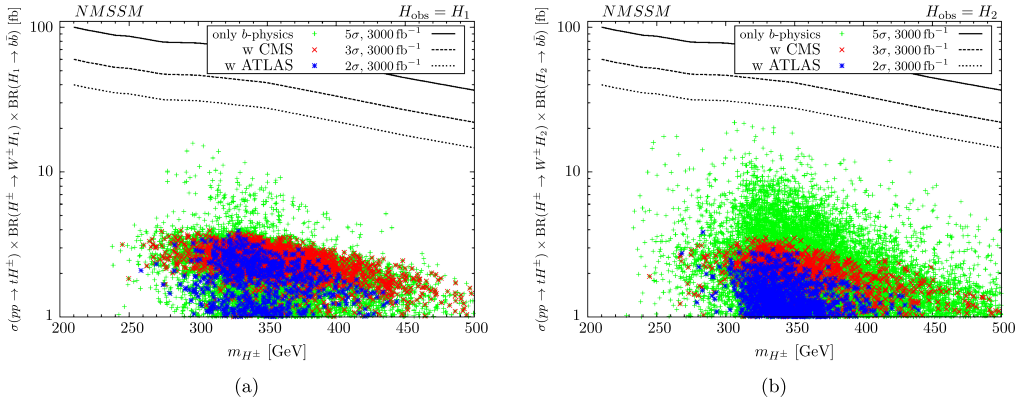


Fig. 5. Signal cross section as a function of  $m_{H^\pm}$  in the NMSSM when (a)  $H_{\text{obs}} = H_1$  and (b)  $H_{\text{obs}} = H_2$ . See text for details. (For interpretation of the colors in this figure, the reader is referred to the web version of this article.)

these constraints further reduces the maximum signal cross section obtainable to below 5 fb. For the case with  $H_{\text{obs}} = H_2$  the signal cross section, shown in Fig. 5(b), can reach slightly higher to around 20 pb, for the green points. This is owing to the somewhat larger  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  obtainable for low  $m_{H^\pm}$  in this case compared to the  $H_{\text{obs}} = H_1$  case. However, again the overall signal cross section is highly diminished for points observing the ATLAS or CMS signal rate constraints. Also shown in Figs. 5(a) and (b) are the  $2\sigma$  (exclusion),  $3\sigma$  (evidence) and  $5\sigma$  (discovery) sensitivity curves for  $3000 \text{ fb}^{-1}$  accumulated luminosity at the LHC 14 TeV run. All the good points from the scans lie well below the lowest ( $2\sigma$ ) curve, implying that none of them has a signal cross section large enough to be testable even at such a high luminosity.

### 6.3. 2HDM Types I and II

The scanned ranges of the parameters in these two models are shown in Table 3. Note that in the 2HDM-II,  $m_{H^\pm} \lesssim 320 \text{ GeV}$  is excluded for all values of  $\tan\beta$  by the constraint on  $\text{BR}(\bar{B} \rightarrow X_s \gamma)$ , while  $\tan\beta \lesssim 1.5$  is ruled out for  $m_{H^\pm}$  up to 500 GeV or so by the  $\Delta B_d$  constraint, according to [57]. We therefore reduced the input range of  $m_{H^\pm}$  instead of imposing these constraints during the scans for this model. The  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  for 2HDM-I



Table 3

Ranges of the input parameters scanned for the 2HDM Types I and II.

Parameter	2HDM-I		2HDM-II	
	$H_{\text{obs}} = h$	$H_{\text{obs}} = H$	$H_{\text{obs}} = h$	$H_{\text{obs}} = H$
$m_h$ (GeV)	123–127	80–115	123–127	80–115
$m_H$ (GeV)	135–500	123–127	135–500	123–127
$m_{H^\pm} = m_A$ (GeV)	135–500		320–500	
$\tan \beta$	1.5–6			
$ \sin(\beta - \alpha) $	0–1			
$m_{12}^2$ (GeV <sup>2</sup> )	$0 - m_A^2 \cos \beta \sin \beta$			

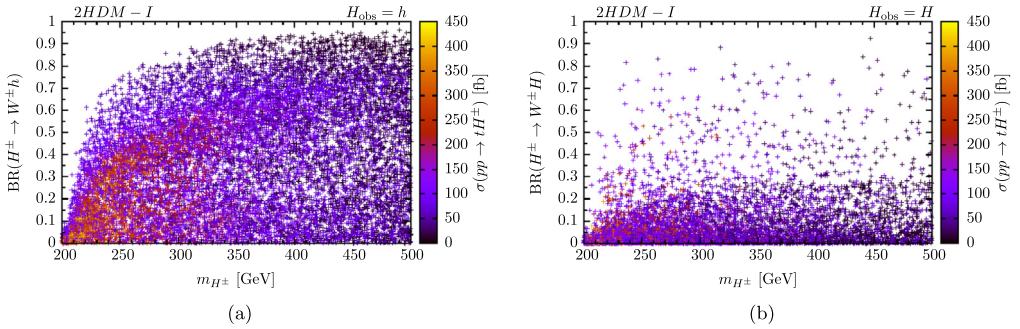


Fig. 6.  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  as a function of  $m_{H^\pm}$  in the 2HDM-I when (a)  $H_{\text{obs}} = h$  and (b)  $H_{\text{obs}} = H$ , with the heat map showing the  $\sigma(pp \rightarrow tH^\pm)$ .

with the  $H_{\text{obs}} = h$  case, shown in Fig. 6(a), can be as high as  $\sim 95\%$  for a fairly large number of points. Moreover, compared to the NMSSM, while the maximum  $\sigma(pp \rightarrow tH^\pm)$  reachable is much lower here, the  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  grows much more sharply with increasing  $m_{H^\pm}$ . As a result, there are plenty of low  $m_{H^\pm}$  points where both the  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  as well as the  $\sigma(pp \rightarrow tH^\pm)$ , shown by the heat map, can be significant. In Fig. 6(b) are shown the corresponding quantities for the  $H_{\text{obs}} = H$  case in the 2HDM-I. In this case a very large  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  is obtainable for a comparatively much smaller number of points and it mostly stays below 40%.

In Fig. 7(a) we show the signal cross section for the  $H_{\text{obs}} = h$  case in the 2HDM-I as a function of  $m_{H^\pm}$ . The colour convention for the points in all the figures showing the signal cross section henceforth is the same as in Fig. 5. We note that, owing to the much larger  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  generally obtainable in this model compared to the NMSSM, the total cross section can reach as high as about 100 fb. A small portion of the green points with  $m_{H^\pm} > 400$  GeV lies above the  $2\sigma$  sensitivity curve corresponding to  $\mathcal{L} = 300 \text{ fb}^{-1}$  and should thus be reachable at the LHC. The picture, however, becomes grim when the LHC signal rate constraints are imposed. Points consistent with the CMS constraints have a maximum possible cross section of around 20 fb, while none of the points obtained in the scans are able to satisfy the ATLAS constraints.

Turning to the 2HDM-II, for the  $H_{\text{obs}} = h$  case one sees in Fig. 8(a) that in this model both the  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  and the  $\sigma(pp \rightarrow tH^\pm)$  show a similar behaviour as noted in the 2HDM-I above, being significantly large simultaneously for a number of points with  $m_{H^\pm}$  up to  $\sim 400$  GeV. The maximum obtainable values of both these quantities are also similar to those in the 2HDM-I. In the  $H_{\text{obs}} = H$  case the  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  struggles to reach high values

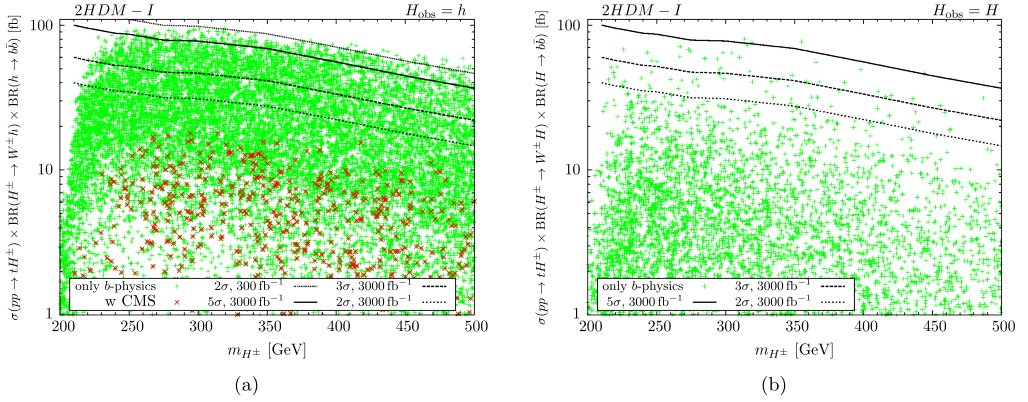


Fig. 7. Signal cross section as a function of  $m_{H^\pm}$  in the 2HDM-I when (a)  $H_{\text{obs}} = h$  and (b)  $H_{\text{obs}} = H$ . See text for details.

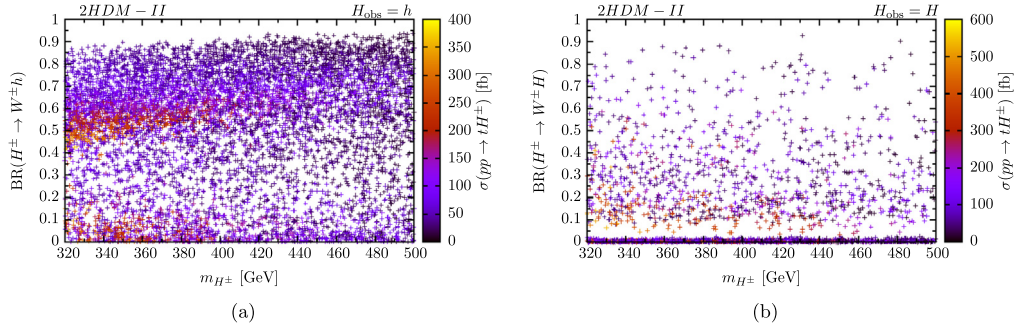


Fig. 8.  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  as a function of  $m_{H^\pm}$  in the 2HDM-II when (a)  $H_{\text{obs}} = h$  and (b)  $H_{\text{obs}} = H$ , with the heat map showing the  $\sigma(pp \rightarrow tH^\pm)$ .

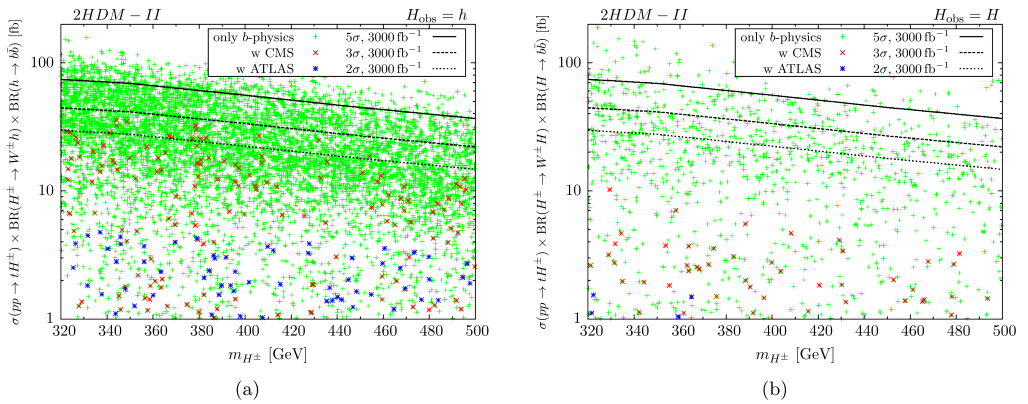


Fig. 9. Signal cross section as a function of  $m_{H^\pm}$  in the 2HDM-II when (a)  $H_{\text{obs}} = h$  and (b)  $H_{\text{obs}} = H$ . See text for details.

Table 4  
Ranges of the input parameters scanned for the A2HDM.

Parameter	$H_{\text{obs}} = h$	$H_{\text{obs}} = H$
$m_h$ (GeV)	123–127	80–115
$m_H$ (GeV)	135–300	123–127
$m_{H^\pm} = m_A$ (GeV)	200–500	
$ \sin \alpha $	0–1	
$\lambda_2$	$0-4\pi$	
$\lambda_3$	$-\sqrt{\lambda_1 \lambda_2} - 4\pi$	
$ \lambda_7 $	$0-4\pi$	
$ \beta^{U,D,L} $	0–1.57	

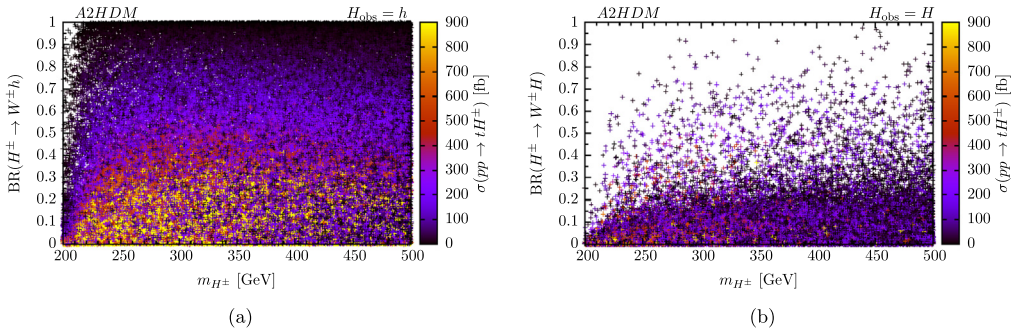


Fig. 10.  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  as a function of  $m_{H^\pm}$  in the A2HDM when (a)  $H_{\text{obs}} = h$  and (b)  $H_{\text{obs}} = H$ , with the heat map showing the  $\sigma(pp \rightarrow tH^\pm)$ .

generally and in fact stays close to 0 for a vast majority of the points, as seen in Fig. 8(b). In Figs. 9(a) and (b) we show the signal cross sections for the  $H_{\text{obs}} = h$  and  $H_{\text{obs}} = H$  cases, respectively, in the 2HDM-II. In the former case, not only do a large number of points observing only the  $b$ -physics constraints lie above the  $5\sigma$  sensitivity curve for  $\mathcal{L} = 3000 \text{ fb}^{-1}$ , but also some of the points consistent with the CMS constraints can have a signal cross section in excess of 30 fb and should thus be accessible at the LHC. In the  $H_{\text{obs}} = H$  case, however, the maximum reachable cross section for points consistent with the CMS and ATLAS signal rate constraints barely exceeds 10 fb and 1.5 fb, respectively, only when  $m_{H^\pm}$  is below 350 GeV or so.

#### 6.4. A2HDM

The scanned ranges of the A2HDM parameters are given in Table 4 and have been adopted from [69]. In Fig. 10(a) we show the  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  for the  $H_{\text{obs}} = h$  case, which can reach unity over the entire desired mass range of  $H^\pm$ . Also, the  $\sigma(pp \rightarrow tH^\pm)$ , illustrated by the heat map in the figure, can reach the pb level, but it is maximal only for points for which the  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  is relatively small,  $\lesssim 40\%$ . On the other hand, Fig. 10(b) shows that in the  $H_{\text{obs}} = H$  case the  $\text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}})$  mostly stays below  $\sim 35\%$ .

In Fig. 11(a) the signal cross section for the  $H_{\text{obs}} = h$  case is shown. This cross section can reach much higher,  $\sim 700 \text{ fb}$ , than in the ordinary 2HDMs, when the constraints from the LHC Higgs boson searches are not imposed. Points with such a high cross section lie above even the  $5\sigma$  sensitivity curve for the LHC with  $\mathcal{L} = 300 \text{ fb}^{-1}$ . This implies that the  $H^\pm$  in this model could be discoverable at the standard luminosity LHC over almost the entire mass range analysed for

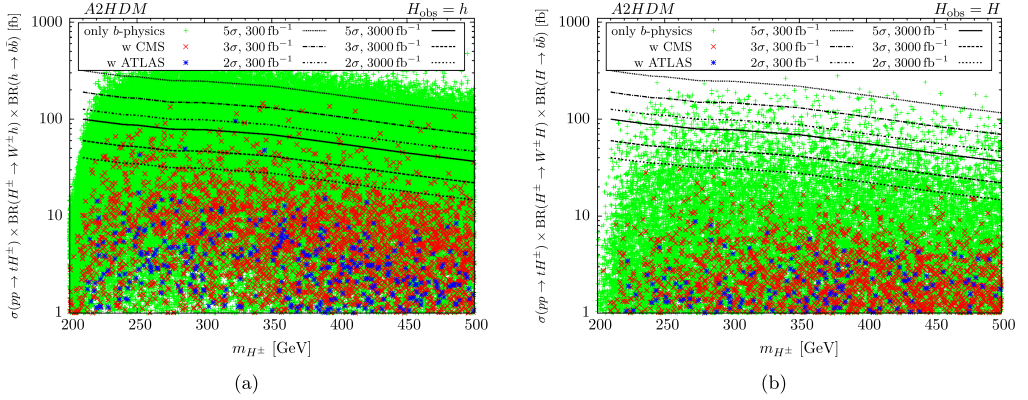


Fig. 11. Signal cross section as a function of  $m_{H^\pm}$  in the A2HDM when (a)  $H_{\text{obs}} = h$  and (b)  $H_{\text{obs}} = H$ . See text for details.

this channel. However, as in the other models above, points satisfying the LHC constraints have a much smaller signal cross section generally. Still, unlike in any of the other models considered here, a small number of points consistent with the CMS constraints lies above the  $5\sigma$  sensitivity curve for  $\mathcal{L} = 3000 \text{ fb}^{-1}$  and could thus be visible at the high luminosity LHC. The same is not true though for the  $H_{\text{obs}} = H$  case, seen in Fig. 11(b), where only a couple of points consistent with the CMS constraints appear to be testable at the high luminosity LHC.

## 7. Conclusions

In this article we have analysed the detectability of  $H^\pm$  in the  $WH_{\text{obs}}$  decay mode in some minimal extensions of the SM, at the upcoming Run 2 of the LHC with  $\sqrt{s} = 14 \text{ TeV}$ . We have discussed some important features of the models of our interest, in particular the coupling parameters governing the production of  $H^\pm$  in  $pp$  collisions as well as the  $H^\pm \rightarrow WH_{\text{obs}}$  decay process. We have performed dedicated scans of the parameter spaces of these models to search for their regions where a  $H^\pm$  with a mass lying in the 200 GeV–500 GeV range can be obtained and its production cross section can be maximised. These scans were subject to the most relevant constraints from  $b$ -physics, from the LHC Higgs boson searches and, in the case of SUSY models, from relic density measurements. Moreover, in the NMSSM as well as in the 2HDMs we considered both the possibilities of the observed Higgs boson being the lightest or the next-to-lightest CP-even scalar of the model.

We then reconstructed the signal and the background in the  $bbb(b)jj\ell\nu_\ell$  final state and, through a dedicated detector-level analysis, estimated the signal significance for various accumulated luminosities at the LHC. We found that, through a judicious choice of selection criteria, including a veto on  $t\bar{t}$  events and the requirement of a reconstructed 125 GeV Higgs boson from a pair of  $b$ -tagged jets, we were able to significantly reduce the backgrounds. The semi-leptonic channel provides enough kinematic information to reconstruct the  $m_{H^\pm}$  peak and identify signals with  $\sigma(pp \rightarrow tH^\pm) \times \text{BR}(H^\pm \rightarrow W^\pm H_{\text{obs}}) \times \text{BR}(H_{\text{obs}} \rightarrow b\bar{b}) \sim \mathcal{O}(100 \text{ fb})$  with an integrated luminosity of  $300 \text{ fb}^{-1}$ , with even better sensitivity at high luminosities.

We have concluded that in the SUSY models studied here, the  $H^\pm \rightarrow WH_{\text{obs}}$  decay channel does not carry as much promise for the identification of a  $H^\pm$  as has been envisaged in some earlier studies. This is due to the fact that the  $pp \rightarrow H^\pm$  production process and the sub-



sequent  $H^\pm \rightarrow WH_{\text{obs}}$  decay process generally show contrasting dependence on the various parameters involved. The situation looks a bit better in the  $Z_2$ -symmetric 2HDMs, as long as the constraints from the LHC measurements of the Higgs boson signal rates are ignored. Imposing these constraints leaves an insignificant number of points in the 2HDM-II visible at only the high luminosity ( $\sim 3000 \text{ fb}^{-1}$ ) LHC, implying that the Higgs boson assumed to be the one observed at the LHC in these scenarios deviates substantially from SM-like properties. In the case of the A2HDM, a fairly large portion of the parameter space could in general be tested even at the standard luminosity ( $\sim 300 \text{ fb}^{-1}$ ) LHC. However, again if the measurements of the observed Higgs boson signal rates do not fluctuate much from the current ones, only a few parameter space points lie within the reach of the LHC at this luminosity.

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