PHYSICAL REVIEW D **83,** 075006 (2011)

$W^{\pm}H^{\mp}$ associated production at the LHC in the general two-Higgs-doublet model with spontaneous *CP* violation

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(Received 15 November 2010; published 11 April 2011)

Spontaneous *CP* violation motivates the introduction of two Higgs doublets in the electroweak theory. Such a simple extension of the standard model has three neutral Higgs bosons and a pair of charged Higgs; it especially leads to rich *CP*-violating sources, including the induced Kobayashi-Maskawa *CP*-violating phase, the mixing of the neutral Higgs bosons due to the *CP*-odd Higgs, and the effective complex Yukawa couplings of the charged and neutral Higgs bosons. Within this model, we present the production of a charged Higgs boson in association with a *W* boson at the LHC, and calculate in detail the cross section and the transverse momentum distribution of the associated *W* boson.

DOI: 10.1103/PhysRevD.83.075006 PACS numbers: 12.60.Fr, 14.80.Bn, 14.80.Fd

I. INTRODUCTION

In the standard model (SM), the fermions and gauge bosons get masses through the Higgs mechanism with a single weak-isospin doublet Higgs field. After the electroweak symmetry breaking, three Goldstone modes are absorbed to build up the longitudinal W and Z gauge bosons, and only one physical scalar called the SM Higgs boson is physical. Since the exact breaking mechanism is not very clear, and the Higgs have not been detected yet, many extension models have been proposed.

One of the simplest extensions of the SM is to add an extra Higgs doublet motivated from spontaneous *CP* violation (SCPV) [1–6]. It has been shown that if one Higgs doublet is needed for the mass generation, then an extra Higgs doublet is necessary for the spontaneous *CP* violation to explain the origin of *CP* violation in SM. In such a model, the *CP* violation is originated from a single relative phase of two vacuum expectation values, which not only gives an explanation for the Kobayashi-Maskawa *CP*-violating mechanism [7] in the SM, but also leads to a new type of *CP*-violating source [4,5]. Such a two-Higgs-doublet model (2HDM) is also called Type III 2HDM to distinguish it from the Types I and II 2HDM.

The common feature of the three types of 2HDM is that after $SU(2)_L \otimes U(1)_Y$ gauge symmetry spontaneous breaking, there are three neutral Higgs and one pair of charged Higgs. As shown in our previous work [8], the neutral Higgs bosons decay to $b\bar{b}$ when their masses are light, which is difficult to be detected due to the strong background at the LHC. Therefore, the charged Higgs boson ($H^{\pm}$) is of special interest, since there are no charged scalars in the SM, and thus its discovery would constitute an indisputable proof of physics beyond the standard model. Thus, the hunt for charged Higgs bosons will play a central role in the search for new physics at the LHC experiments.

Currently, the limit or constraint to the charged Higgs mass is not very strong and is also model-dependent. The best model-independent direct limit from the LEP experiments is $m_{H^{\pm}} > 78.6$ GeV (95% CL) [9], assuming only the decays $H^+ \to c\bar{s}$ and $H^+ \to \tau\nu_{\tau}$. And, as the charged Higgs will contribute to flavor-changing neutral currents (FCNC) at one-loop level, the indirect constraint can be extracted from B-meson decays. In the Type II model, the constraint is $M_{H^{\pm}} \gtrsim 350$ GeV for $\tan\beta$ larger than 1, and even stronger for smaller $\tan\beta$ [10]. However, as the phases of the Yukawa couplings in Type III model are free, $m_{H^{\pm}}$ can be as low as 100 GeV [11]. In this note, we take it as free from 150 GeV to 500 GeV.

At the LHC, the interesting channels for the charged boson production are $gb \to H^- t$ for $m_{H^{\pm}} > m_t + m_b$ and $gg \to H^- t\bar{b}$ for $m_{H^{\pm}} \lesssim m_t - m_b$ [12–15]. In these channels, the leptonic decay $H^+ \to \tau^+\nu$ seems most promising for detecting light charged Higgs, while the hadronic decay $H^+ \to t\bar{b}$ may be hopeful above a threshold with efficient b-tagging [16–23]. Another interesting channel is to produce the $H^{\pm}$ in association with W bosons, and the leptonic decays of the W-boson can serve as a trigger for the $H^{\pm}$ boson search. This channel can also cover the transition region search, $M_{H^{\pm}} \sim m_t$. The dominant channels for $W^{\mp}H^{\pm}$ production are $b\bar{b} \to W^{\pm}H^{\mp}$ at tree level and $gg \to W^{\pm}H^{\mp}$ at one-loop level. $W^{\mp}H^{\pm}$ production at hadron colliders in Type II 2HDM and the minimal supersymmetric standard model has been extensively studied in [24–30]. The *CP* violation effect is also explored at the muon collider [31]. In this paper, we will study it in the Type III 2HDM with Spontaneous *CP* Violation (SCPV) [4,5]. The discovery of relative light charged Higgs boson with $M_{H^{\pm}} <$ 350 GeV distinguishes it from the Type II 2HDM.

This paper is organized as follows. In Sec. II, we shall first give a brief introduction of the 2HDM with SCPV and

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some conventions. Then, calculations are outlined in Sec. III and numerical results are shown in Sec. IV. Finally, we come to our conclusions.

II. 2HDM WITH SCPV

We begin with a brief introduction to the model by showing the spontaneous *CP* violation and its difference to Type I and Type II models. The two complex Higgs doublets are generally expressed as

$$\Phi_1 = \begin{pmatrix} \phi_1^+ \\ \phi_1^0 \end{pmatrix}, \qquad \Phi_2 = \begin{pmatrix} \phi_2^+ \\ \phi_2^0 \end{pmatrix}, \tag{1}$$

and the potential is

$$\begin{aligned} V(\Phi_1,\Phi_2) = & -\mu_1^2\Phi_1^\dagger\Phi_1 - \mu_2^2\Phi_2^\dagger\Phi_2 - (\mu_{12}^2\Phi_1^\dagger\Phi_2 + \mathrm{H.c.}) \\ & + \lambda_1(\Phi_1^\dagger\Phi_1)^2 + \lambda_2(\Phi_2^\dagger\Phi_2)^2 + \lambda_3(\Phi_1^\dagger\Phi_1\Phi_2^\dagger\Phi_2) \\ & + \lambda_4(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1) + \frac{1}{2}[\lambda_5(\Phi_1^\dagger\Phi_2)^2 + \mathrm{H.c.}] \\ & + [(\lambda_6\Phi_1^\dagger\Phi_1 + \lambda_7\Phi_2^\dagger\Phi_2)(\Phi_1^\dagger\Phi_2) + \mathrm{H.c}]. \end{aligned} \tag{2}$$

With λ_5 nonzero and real, *CP* violation can arise from nonzero values of one or more of μ_{12}^2, λ_6 or λ_7. If these three (and λ_5) are all real, *CP* violation can occur spontaneously [1] when $\lambda_5 > 0$, because of the relative phase between the vacuum expectation values. The most interesting case is that only the dimension-2 term μ_{12}^2 is complex, which is known as a soft *CP*-violating phase. Then, it can easily be demonstrated that once all other couplings in the Higgs potential are required to be positive real, the relative *CP*-violating phase of two vacuum expectation values is solely determined by the explicit soft *CP*-violating phase [5,6] via the minimal conditions of Higgs potential. In this case, the *CP* violation remains originating from the spontaneous breaking of symmetry in vacuum, while it can avoid the so-called domain-wall problem.

The Yukawa interaction terms have the following general form:

$$-\mathcal{L}_Y = \eta_{ij}^{(k)}\bar{\psi}_L^i\tilde{\Phi}_k U_R^j + \xi_{ij}^{(k)}\bar{\psi}_L^i\Phi_k D_R^j + \mathrm{H.c}, \tag{3}$$

where $\psi_L^i = (U_L^i, D_L^i)^T$, $\tilde{\Phi}_k = i\tau_2\Phi_k^*$, $\eta_{ij}^{(k)}$ and $\xi_{ij}^{(k)}$ are all real Yukawa coupling constants to keep the interactions *CP*-invariant. The above interactions may lead to flavor-changing neutral currents (FCNC) at the tree level through the neutral Higgs boson exchanges as the Yukawa matrices may not be diagonal. FCNC processes should be strongly suppressed based on the experimental observations. Usually, an *ad hoc* discrete symmetry [32] is often imposed:

$$\begin{aligned} &\Phi_1 \to -\Phi_1 \quad \text{and} \quad \Phi_2 \to \Phi_2, \\ &U_R \to -U_R \quad \text{and} \quad D_R \to \mp D_R, \end{aligned} \tag{4}$$

which correspond to Type I and Type II 2HDM, relying on whether the up- and down-type quarks are coupled to the same or different Higgs doublet. Some interesting phenomena for various cases in such types of models without FCNC have been investigated in detail in [33,34]. When the discrete symmetry is introduced in the potential Eq. (2), it leads to $\mu_{12} = 0$ and $\lambda_6 = \lambda_7 = 0$ and then no spontaneous *CP*-violation any more [35]. Since the FCNC is observed in experiments in weak interactions though it is strongly suppressed, we shall abandon the discrete symmetry and consider the small off-diagonal Yukawa couplings. The naturalness for such small Yukawa couplings may be understood from the approximate global $U(1)$ family symmetries [4,5,36,37]. As if all the up-type quarks and also the down-type quarks have the same masses and no mixing, the theory has an $U(3)$ family symmetry for three generations, while when all quarks have different masses but still no mixing, the theory has the $U(1) \otimes U(1) \otimes U(1)$ family symmetries and the Cabibbo-Kobayashi-Maskawa quark-mixing matrix is a unit matrix. In this case, both the direct FCNC and induced FCNC are absent. In the real world, there are some FCNC processes observed, thus the $U(1)$ family symmetries should be broken down. As all the observed FCNC processes are strongly suppressed, the theory should possess approximate $U(1)$ family symmetries with small off-diagonal mixing among the generations. In this sense, the approximate $U(1)$ family symmetries are enough to ensure the naturalness of the observed smallness of FCNC.

As in the potential Eq. (2), the neutral Higgs bosons will get the vacuum expectation values as follows:

$$\langle\phi_1^0\rangle = \frac{1}{\sqrt{2}}v_1 e^{i\delta_1}, \qquad \langle\phi_2^0\rangle = \frac{1}{\sqrt{2}}v_2 e^{i\delta_2}, \tag{5}$$

where one of the phases can be rotated away due to the global $U(1)$ symmetry in the potential and Yukawa terms. Without losing generality, we may take $\delta_1 = 0$ and $\delta_2 = \delta$. It is then convenient to make a unitary transformation

$$\begin{pmatrix} H_1 \\ H_2 \end{pmatrix} = U\begin{pmatrix} \Phi_1 \\ \Phi_2 \end{pmatrix}, \text{ with } U = \begin{pmatrix} \cos\beta & \sin\beta e^{-i\delta} \\ -\sin\beta & \cos\beta e^{-i\delta} \end{pmatrix}, \tag{6}$$

where $\tan\beta = v_2/v_1$. After making the above transformation and redefining the ϕ_i^0, we can reexpress the Higgs doublets as follows:

$$H_1 = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v + \phi_1^0 \end{pmatrix} + \mathcal{G}, \quad H_2 = \frac{1}{\sqrt{2}}\begin{pmatrix} \sqrt{2}H^+ \\ \phi_2^0 + i\phi_3^0 \end{pmatrix}, \tag{7}$$

with $v^2 = v_1^2 + v_2^2$ and $v \simeq 246$ GeV, which is the same as in the standard model. Thus, in this new basis, only the Higgs doublet H_1 gives masses to the gauge bosons, quarks, and leptons. The Higgs fields $\mathcal{G}$ are the goldstone particles absorbed by the gauge bosons, while $H^\pm$ are mass eigenstates of the charged scalar Higgs. $\phi_i^0 = (\phi_1^0, \phi_2^0, \phi_3^0)$ are the neutral Higgs bosons in the electroweak eigenstates. In general, they are not the same as the physics Higgs bosons $h_j = (h_1, h_2, h_3)$ in the mass eigenstates, but related via an orthogonal SO(3) transformation

$$\phi_i^0 = O_{ij}h_j \quad \text{with} \quad i, j = 1, 2, 3, \tag{8}$$

where O_{ij} depends on the λ_i and μ_i in the Higgs potential. When there is no mixing between ϕ_1^0, ϕ_2^0, and ϕ_3^0, h_1, h_2, and h_3 then correspond to h^0, H^0 (CP-even) and A^0 (CP-odd) in the literature, respectively. For the convenience of later discussion, we will denote the mixing angle α of h_1 and h_3 by meaning that

$$\begin{pmatrix} \phi_1^0 \\ \phi_2^0 \\ \phi_3^0 \end{pmatrix} = \begin{pmatrix} \cos\alpha & 0 & -\sin\alpha \\ 0 & 1 & 0 \\ \sin\alpha & 0 & \cos\alpha \end{pmatrix} \begin{pmatrix} h_1 \\ h_2 \\ h_3 \end{pmatrix}. \tag{9}$$

In the new basis of Eq. (7), the Yukawa interaction terms in Eq. (3) can be reexpressed as

$$-\mathcal{L}_Y = \eta_{ij}^U \bar{\psi}_L^i \tilde{H}_1 U_R^j + \xi_{ij}^U \bar{\psi}_L^i \tilde{H}_2 U_R^j + \eta_{ij}^D \bar{\psi}_L^i H_1 D_R^j + \xi_{ij}^D \bar{\psi}_L^i H_2 D_R^j + \text{H.c}, \tag{10}$$

where

$$\begin{aligned} \eta_{ij}^U &= \eta_{ij}^{(1)} \cos\beta + \eta_{ij}^{(2)} e^{-i\delta} \sin\beta \equiv \sqrt{2} M_{ij}^U / v, \\ \xi_{ij}^U &= -\eta_{ij}^{(1)} \sin\beta + \eta_{ij}^{(2)} e^{-i\delta} \cos\beta, \\ \eta_{ij}^D &= \xi_{ij}^{(1)} \cos\beta + \xi_{ij}^{(2)} e^{i\delta} \sin\beta \equiv \sqrt{2} M_{ij}^D / v, \\ \xi_{ij}^D &= -\xi_{ij}^{(1)} \sin\beta + \xi_{ij}^{(2)} e^{i\delta} \cos\beta \end{aligned} \tag{11}$$

and $M^{U,D}$ are fermion mass matrices. As the Yukawa coupling terms η^U and ξ^D become complex due to the vacuum phase δ with real $\eta^{(1),(2)}$ and $\xi^{(1),(2)}$ defined in Eq. (3), the resulting mass matrices are also complex. Then unitary transformations are needed to diagonalize the mass matrices,

$$\begin{aligned} u_{L,R}^j &= V_{L,R}^{jk,U} U_{L,R}^k, V_L^U \eta^U V_R^{U\dagger} = \sqrt{2} \frac{m_u}{v}; \\ d_{L,R}^j &= V_{L,R}^{jk,D} D_{L,R}^k, \qquad V_L^D \eta^D V_R^{D\dagger} = \sqrt{2} \frac{m_d}{v}. \end{aligned} \tag{12}$$

By transforming electroweak interaction eigenstates of the fermions, $U_{L,R}$ and $D_{L,R}$, to their mass eigenstates, $u_{L,R}$ and $d_{L,R}$, we denote the final Yukawa couplings in the quark mass eigenstates by $\xi^{u,d}$, with the relation,

$$\xi^{u,d} = V_L^{U,D} \xi^{U,D} V_R^{U,D\dagger}. \tag{13}$$

With the above notation, the Yukawa interaction terms are

$$\begin{aligned} -\mathcal{L}_Y = \sum_{i=1}^{3} &\left[m_u^i \bar{u}_L^i u_R^i \left(1 + \frac{\phi_1^0}{v}\right) + m_d^i \bar{d}_L^i d_R^i \left(1 + \frac{\phi_1^0}{v}\right) \right] \\ &+ \frac{1}{\sqrt{2}} \bar{u}_L^i \xi_{ij}^u u_R^j (\phi_2^0 - i\phi_3^0) + \frac{1}{\sqrt{2}} \bar{d}_L^i \xi_{ij}^d d_R^j (\phi_2^0 + i\phi_3^0) \\ &- \bar{d}_L^i \hat{\xi}_{ij}^u u_R^j H^- + \bar{u}_L^i \hat{\xi}_{ij}^d d_R^j H^+ + \text{H.c} \end{aligned} \tag{14}$$

where the charged Yukawa coupling are defined as

$$\hat{\xi}^u = V_{\text{CKM}}^\dagger \xi^u; \qquad \hat{\xi}^d = V_{\text{CKM}} \xi^d. \tag{15}$$

It can be seen that when there was no mixing among ϕ_1^0, ϕ_2^0, and ϕ_3^0, the scalar ϕ_1^0 plays the role of the Higgs in the SM.

In the following discussions, we shall use ξ^u, ξ^d, and the quarks' masses as the free and independent input parameters instead of the original Yukawa coupling matrices ($\eta_{ij}^{(k)}$, $\xi_{ij}^{(k)}$) given in Eq. (3) and the parameter β. It is convenient to parameterize the Yukawa couplings $\xi_{ij}^{u,d}$ by using the quark mass scales [38],

$$\xi_{ij}^{u,d} \equiv \lambda_{ij} \sqrt{2 m_i^{u,d} m_j^{u,d}} / v, \tag{16}$$

where the i, j are the flavor indexes (for ξ_{ij}^u, i, $j = u$, c, t and for ξ_{ij}^d, i, $j = d$, s, b), and the smallness of the off-diagonal elements are characterized by the hierarchical mass scales of quarks and the parameters λ_{ij}.

After the transformation of the Higgs in Eq. (6), the gauge part of the Higgs in the basis of ϕ_i^0 can be written as [4]

$$\begin{aligned} \mathcal{L}_{\text{G}} &= (D_\mu H_1)^\dagger (D^\mu H_1) + (D_\mu H_2)^\dagger (D^\mu H_2) \\ &= \frac{1}{2} \partial_\mu \phi_1^0 \partial^\mu \phi_1^0 + \frac{(v + \phi_1^0)^2}{8} [(g'^2 + g^2) Z^2 + 2g^2 W^+ W^-] + \frac{1}{2} (\partial \phi_2^0 \partial^\mu \phi_2^0 + \partial \phi_3^0 \partial^\mu \phi_3^0) + \partial_\mu H^- \partial^\mu H^+ \\ &\quad + e^2 H^+ H^- A^2 + \frac{(g^2 - g'^2)^2}{4(g^2 + g'^2)} H^+ H^- Z^2 + \frac{g^2}{2} H^+ H^- W^+ W^- + \frac{e(g^2 - g'^2)}{\sqrt{g^2 + g'^2}} H^+ H^- Z \cdot A + \frac{g^2}{4} W^+ W^- (\phi_2^{02} + \phi_3^{02}) \\ &\quad + \frac{g^2 + g'^2}{8} (\phi_2^{02} + \phi_3^{02}) Z^2 + \Big\{ ie A^\mu H^- \partial_\mu H^+ + i \frac{g^2 - g'^2}{2\sqrt{g^2 + g'^2}} Z^\mu H^- \partial_\mu H^+ + \frac{ig}{2} W_\mu^- (\phi_2^0 - i\phi_3^0) \partial^\mu H^+ \\ &\quad + \frac{eg}{2} A^\mu W_\mu^- H^+ (\phi_2^0 - i\phi_3^0) + \frac{g}{4} \frac{g^2 - g'^2}{\sqrt{g^2 + g'^2}} H^+ W_\mu^- Z^\mu (\phi_2^0 - i\phi_3^0) + \frac{ig}{2} H^- W_\mu^+ (\partial^\mu \phi_2^0 + i\partial^\mu \phi_3^0) \\ &\quad + \frac{i\sqrt{g^2 + g'^2}}{4} (\phi_2^0 - i\phi_3^0) Z^\mu (\partial_\mu \phi_2^0 + i\partial_\mu \phi_3^0) - \frac{g\sqrt{g^2 + g'^2}}{4} H^- W_\mu^+ Z^\mu (\phi_2^0 + i\phi_3^0) + \text{H.c.} \Big\}. \end{aligned} \tag{17}$$

TABLE I. The decay width of the Higgs h_3 when $m_{h_3} =$ 500 GeV. We list the fractional width $W^\pm H^\mp$ and the sum width of $b\bar{b}$, $t\bar{t}$, $WW^{(*)}$, and $ZZ^{(*)}$ separately.

Γ_{h_3}(GeV)	Case A	Case B	Case C	$W^\mp H^\pm$(200 GeV)
$\alpha = 0$	17.552	43.961	62.626	44.094
$\alpha = \pi/4$	41.764	54.969	64.301	22.047
$\alpha = \pi/2$	65.135	65.135	65.135	0

III. $W^\pm H^\mp$ ASSOCIATED PRODUCTION

At hadron colliders, the dominant mechanisms for $W^\pm H^\mp$ associated production are $b\bar{b}$ annihilation at tree level and gluon-gluon fusion at one-loop level. As the Feynman diagrams show in Fig. (1), the $b\bar{b}$ annihilation proceeds either via s-channel resonance mediated by the neutral Higgs $h_i s$, or by t-channel dominated by the top quark exchange. Here, we treat b and $\bar{b}$ quarks as active partons inside the colliding protons and use the parton distribution functions set [39].

An alternative $W^\pm H^\mp$ production mechanism is provided by gluon-gluon fusion shown in Fig. (2). From the Feynman diagrams, although the parton-level cross section of gluon fusion is suppressed by α_s^2 relative to the one of $b\bar{b}$ annihilation, it is expected to yield a comparable contribution at 14 TeV hadron colliders, due to the overwhelming gluon luminosity. But as our result, shown in Figs. (3 and 4), the $\sigma(gg \to H^\pm W^\mp)$ is much smaller than $\sigma(b\bar{b} \to H^\pm W^\mp)$, unless the Yukawa of top quark ξ_t is very large.

The relevant interaction terms in this calculation are

$$\mathcal{L}_{h_j H^\pm W^\mp} = \frac{g}{2}\sum_j [\tilde{g}_j (h_j i \overset{\partial_\mu}{\leftrightarrow} H^-) W^{+,\mu} - \tilde{g}_j^* (h_j i \overset{\partial_\mu}{\leftrightarrow} H^+) W^{-,\mu}]$$
$$\tilde{g}_j = O_{2j} + iO_{3j} = g_{h_j H^- W^+}, \tag{18}$$

Yukawa terms (the light quark parts are ingored)

$$\mathcal{L}_{h_i\bar{q}q} = -\frac{g m_q}{2m_W}\sum_i h_i \bar{q}[g_i^q P_L + g_i^{q*} P_R]q,$$
$$q = b \text{ or } t g_i^b = O_{1i} + \lambda_b^*(O_{2i} + iO_{3i}),$$
$$\lambda_b \equiv \lambda_{bb} g_i^t = O_{1i} + \lambda_t^*(O_{2i} - iO_{3i}), \quad \lambda_t \equiv \lambda_{tt}, \tag{19}$$

where $P_L = (1 - \gamma_5)/2$, $P_R = (1 + \gamma_5)/2$, and

$$\mathcal{L}_{H^\mp tb} = \frac{g}{\sqrt{2}m_W} H^+ \bar{t}[\lambda_t^* m_t P_L - \lambda_b m_b P_R]b + \text{H.c.} \tag{20}$$

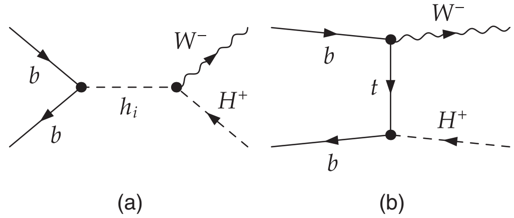

FIG. 1. Feynman diagrams for $W^- H^+$ production via $b\bar{b}$ annihilation at tree level, where h_i denote (h_1, h_2, h_3).

As in the s-channel diagram of $b\bar{b} \to H^\pm W^\mp$, generally when we consider the decay width of Higgs particles in the Higgs propagators

$$S_{h_i} = \frac{1}{p^2 - M_i^2 + iM_i\Gamma_{h_i}} = \frac{p^2 - M_i^2 - iM_i\Gamma_{h_i}}{(p^2 - M_i^2)^2 + M_i^2\Gamma_{h_i}^2}, \tag{21}$$

which is the same for the CP-conjugate processes. Thus, the resulting effective phase of complex production amplitude caused by the Higgs propagator with considering decay width will play the role of strong phase in the hadronic decays. This is different from the phase caused either from the mixing between CP-even and CP-odd Higgs states or from the complex couplings of the electroweak Higgs eigenstates in Eq. (14), which has an opposite sign between the CP-conjugate processes. In this case, there will be CP asymmetry in $W^\pm H^\mp$ productions such as considered in [29,31]. But at the LHC, as the central energy is very high, $s = (p_1 + p_2)^2 \sim \text{TeV}^2$, and the decay width of the Higgs is small, the CP asymmetry is suppressed by the factor Γ_i^2/s, which is hard to be detected. Since the CP violation comes from the interaction terms between different neutral Higgs contributions, to get large CP violation there must exist more than two heavy and very unstable neutral Higgs bosons. On the other hand, as the $pp \to H^\pm W^\pm$ production is dominated by $b\bar{b} \to H^\pm W^\pm$, although the CP asymmetry in $gg \to H^\pm W^\pm$ is larger, the total CP asymmetry of $H^\pm W^\mp$ production on proton-proton collision remains small.

Furthermore, when the three Higgs bosons are all light, we can consider the first order in Eq. (21),

$$S_{h_i} \sim \frac{1}{s} + \mathcal{O}(M_i^2/s), \tag{22}$$

so that the three Higgs bosons have the similar propagators, which makes the effect of the mixing between h_i and h_j be very small, and suppressed by the factor $M_i^2 M_j^2/s^2$ due to the orthogonality of the mixing matrix.

In this paper, we first take the neutral Higgs mixing matrix to be diagonal so as to see the effect of the CP phases of the Yukawa couplings, which are absent in the Type II 2HDM. Then, we consider the effects of the mixing between the Higgs bosons and the dependence of the production on the Higgs masses. In our calculations, the Feynman graphs are generated by using FeynArts [40] and evaluated with FormCalc and LoopTools [41].

IV. NUMERICAL RESULTS

It was observed in the Type III 2HDM [4,5] that the charged Higgs interactions involving the Yukawa couplings $\hat{\xi}^{u,d}$ in Eq. (14) lead to a new type of CP-violating FCNC, even if the neutral current couplings $\xi^{u,d}$ are diagonal. For the parameters concerning the third generation, we may express as

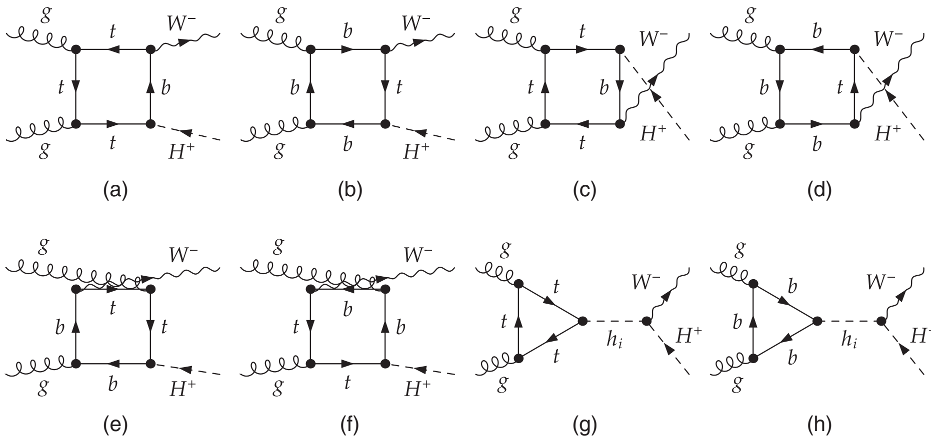

FIG. 2. Feynman diagrams for W^-H^+ production via gg fusion at one-loop level.

$$\xi^u_{tt} = \xi^t = \xi_t e^{\delta_t}, \quad \xi^d_{bb} = \xi^b = \xi_b e^{\delta_b}. \tag{23}$$

The general constraints on the FCNC and the relevant parameter spaces have been investigated in [11,42–46].

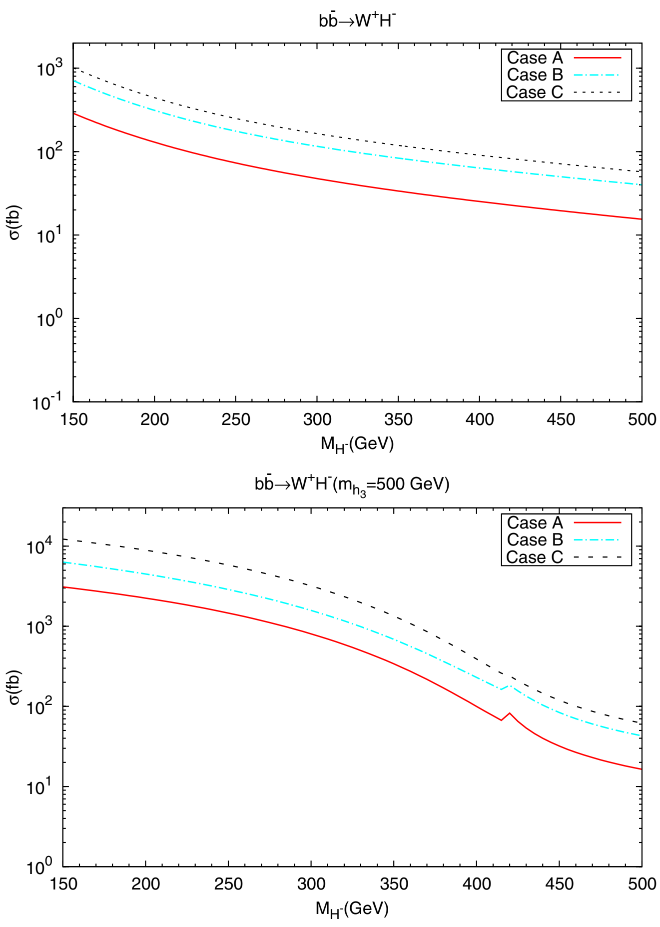

FIG. 3 (color online). The total cross section of W^+H^- production from $b\bar{b}$ channel at the LHC ($\sqrt{S} = 14$ TeV) as function of the charged Higgs's mass in three cases: Case A (solid line), Case B (dashed-dot), and Case C (dashed) with $m_{h_3} = 120$ GeV, 500 GeV.

Here, we may consider the following three typical parameter spaces for the neutral Yukawa couplings of b-quark and t-quark $\xi^q/\sqrt{2} = \lambda_q m_q/v$:

$$\begin{aligned} &\text{Case A: } |\xi^t/\sqrt{2}| = 0.2(\lambda_t = 0.3);\\ &|\xi^b/\sqrt{2}| = 0.5(\lambda_b = 30),\\ &\text{Case B: } |\xi^t/\sqrt{2}| = 0.1(\lambda_t = 0.15);\\ &|\xi^b/\sqrt{2}| = 0.8(\lambda_b = 50), \text{Case C: } |\xi^t/\sqrt{2}|\\ &\qquad = 0.01(\lambda_t = 0.015);\\ &|\xi^b/\sqrt{2}| = 1.0(\lambda_b = 60), \end{aligned} \tag{24}$$

which is consistent with the current experimental constraints in the flavor sector including the B meson decays , even when the neutral Higgs masses are light. In this note, we take

$$\begin{aligned} m_{h_1} &= 115\ \text{GeV}, \qquad m_{h_2} = 160\ \text{GeV},\\ m_{h_3} &= 120\ \text{GeV} \quad \text{or} \quad 500\ \text{GeV} \end{aligned} \tag{25}$$

as input. The strong coupling constant $\alpha_s(\mu)$ is running with $\alpha_s(M_Z) = 0.1176$ [49]. We identify the renormalization and factorization scales with the $W^{\pm}H^{\mp}$ invariant mass. For the charged Higgs mass, the direct limit from LEP is $m_{H^{\pm}} > 78.6$ GeV , and can be as low as 100 GeV from B-meson decay [11]. The strong constraints may arise from the radiative bottom quark decay $b \longrightarrow s\gamma$. In fact, its mass was found to be severely constrained from the $b \longrightarrow s\gamma$ decay in the Type II 2HDM; the lower bound on the charged Higgs mass can be as large as $m_{H^+} \simeq 350$ GeV, which corresponds to the special case in the Type III 2HDM with the parameter $|\xi^t||\xi^b| \sim 0.02$ (or $|\lambda_t\lambda_b| \sim 1$) and a relative phase $\delta_t - \delta_b = 180°$. In this note, we take the charged Higgs mass as free from 150 GeV to 500 GeV and take $m_{H^{\pm}} = 200$ GeV as a special case to study the P_T distributions.

In Figs. (3 and 4), we show the dependence of the cross section on the mass of the charged and neutral Higgs.

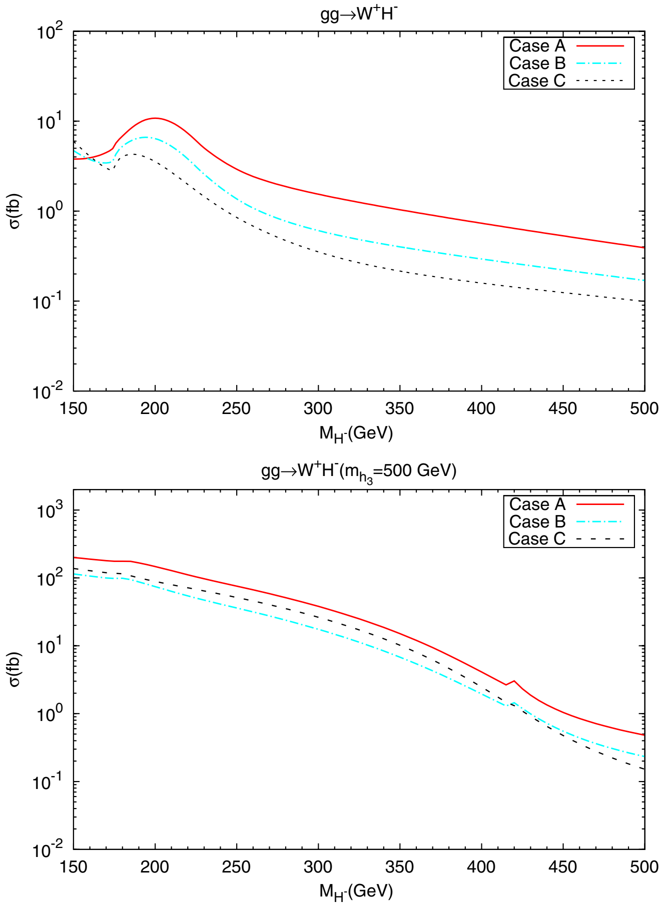

FIG. 4 (color online). The total cross section of W^+H^- production from gg channel at the LHC ($\sqrt{S} = 14$ TeV) as function of the charged Higgs's mass in three cases: Case A (solid line), Case B (dashed-dot), and Case C (dashed) with $m_{h_3} = 120$GeV, 500 GeV.

Compared to the contribution from $b\bar{b}$ with gg, we can see that $b\bar{b}$-annihilation is the dominant channel and its contribution is generally two order larger in Case B and C than that of gg-fusion. In Case A, it is about one order larger. We will only show the diagram from $b\bar{b}$ channel. The $b\bar{b} \to H^{\pm}W^{\mp}$ is dominant by the coupling ξ_b, therefore the $\sigma(b\bar{b} \to H^{\pm}W^{\mp})$ of Case C is the largest and the Case A is the smallest. These results are also confirmed in Type II 2HDM. As the $\tan\beta$ is smaller, the $\sigma(b\bar{b} \to H^{\pm}W^{\mp})$ is smaller while the $\sigma(gg \to H^{\pm}W^{\mp})$ is larger [25]. Now, we focus on the $b\bar{b} \to W^{\mp}H^{\pm}$. We can see that when the neutral Higgs h_3 is heavier, the cross section is larger. That is because when h_3 is heavier, the propagator $S_{h_3} = 1/(s - m_{h_3}^2 + im_{h_3}\Gamma_{h_3})$ is larger and on-shell h_3 can be produced. The shapes of the curves are also changed because of the effect of the width Γ_{h_3}, which plays an important role when $m_{h_3} = 500$ GeV. At about $m_{H^{\mp}} \sim 420$ GeV when $m_{h_3} = 500$ GeV, there is a peak due to the threshold effect.

We also show the differential cross section on the transverse momentum p_T of H^- in Fig. (5) for $b\bar{b}$-channel. As the contribution from gg fusion is small, we do not consider its P_T distribution here and the final result of $pp \to W^{\pm}H^{\mp}$ is dominated by $b\bar{b}$ contribution. The curves have different shapes in the two cases, $m_{h_3} = 500$ GeV and $m_{h_3} = 120$ GeV.

As, in general, the λ_{ij} can be complex, we shall consider the dependence of the cross section on the phase difference between λ_{bb} and λ_{tt} ($\delta = \delta_b - \delta_t$). Generally, we can take $\delta_b = \delta$ and $\delta_t = 0$. If $\delta_t \neq 0$, the curve will be globally shifted, and its shape will not be changed. The cross section of $b\bar{b} \to H^{\pm}W^{\mp}$ channel varies less than 1% as $\delta \in [0, 2\pi]$, as the s- and t-channel are all dominant by ξ_b, and the cross terms are suppressed as $O(\frac{m_b}{m_t})$. Since the total cross section from gg is much smaller than that from $b\bar{b}$ channel in Case B and Case C, λ_{bb} phase has almost negligible effect on the total production of $W^{\mp}H^{\pm}$.

Now, we would like to discuss the effect of the mixing between neutral Higgs bosons. As discussed in last section, the h_1 and h_2 are light; their mixing effect can be neglected. Therefore, for $m_{h_3} = 120$ GeV, the mixing effect

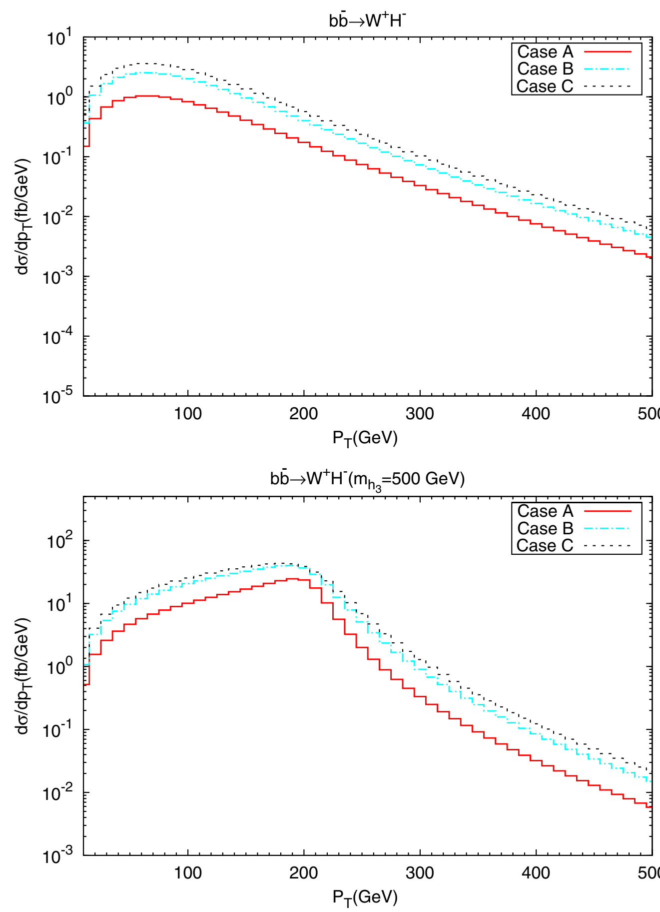

FIG. 5 (color online). The P_T distribution of W^+H^- production from $b\bar{b}$ channel at the LHC ($\sqrt{S} = 14$ TeV) with $m_{h_3} = 120$GeV, 500 GeV and fixed $m_{H^{\pm}} = 200$ GeV in three cases: Case A (solid), Case B (dashed-dot), and Case C (dashed).

to $pp \to W^{\mp}H^{\mp}$ can not be detected. When $m_{h_3} = 500$ GeV, the mixing of h_1 or h_2 with h_3 can be sizable. Here, we shall consider the mixing between h_1 and h_3, as it can be seen from the lagrangian that without mixing, h_1 has no contribution to $W^{\mp}H^{\mp}$ production. The width effect of Higgs bosons will also be included. The total decay widths of h_3 are listed in Table I for different cases and mixing angles α between h_1 and h_3. For simplicity, we only consider the dominant decay channel at tree level $h_3 \to b\bar{b}$, $h_3 \to t\bar{t}$, $h_3 \to WW^{(*)}$, $h_3 \to ZZ^{(*)}$ and $h_3 \to W^{\mp}H^{\pm}$. The $W^{(*)}$ and $Z^{(*)}$ mean that the bosons can be on-shell or off-shell as treated in [50]. As the $b\bar{b}$, $t\bar{t}$, $WW^{(*)}$, and the $ZZ^{(*)}$ modes are not dependent on $m_{H^{\pm}}$, we sum them together and list them with three cases. While the decay mode $h_3 \to W^{\mp}H^{\pm}$ is dependent on the $m_{H^{\mp}}$, and independent of the different cases, we list it separately in the last column only when $m_{H^{\mp}} = 200$ GeV. With the above decay width, we present our results in Fig. (6) for $b\bar{b}$ channel with $m_{h_3} = 500$ GeV. It can be seen that the effect of the mixing for the three cases is similar: as the mixing angle α goes larger, the contribution from the h_3 gets smaller, and the total cross section goes down. And when $\alpha = \pi/2$, the dashed lines are almost the same as the three lines shown in Figs. (3 and 4) with $m_{h_3} = 120$ GeV. This is because in Fig. (3) and Fig. (4), there are only h_2 and h_3 ($m_{h_3} = 120$ GeV) to contribute to the production, and in Fig. (6) with $\alpha = \pi/2$, the roles played by h_3 and h_1 are just interchanged, and the numerical results are also almost the same as $m_{h_1} = 115$ GeV $\sim$ 120 GeV.

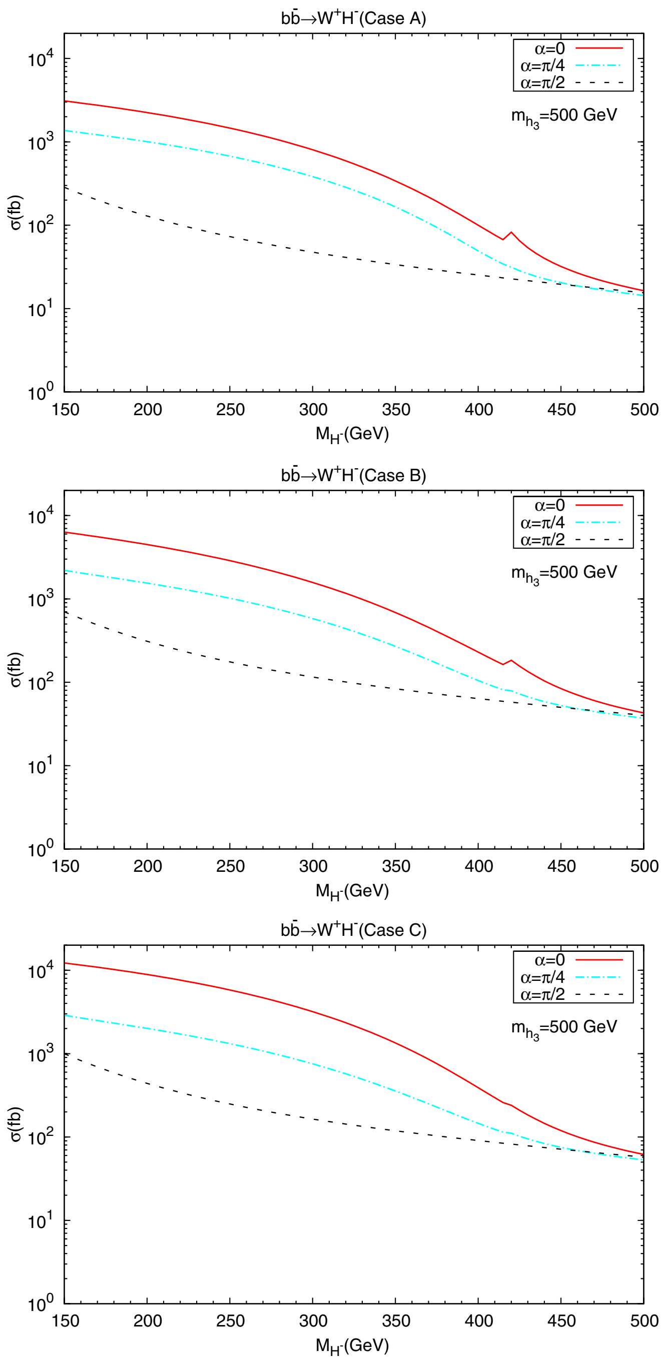

FIG. 6 (color online). The total cross section of W^+H^- production from $b\bar{b}$ channel at the LHC ($\sqrt{S} = 14$ TeV) as the charged Higgs's mass and mixing angle (α) between h_1 and h_3 in three cases: $\alpha = 0$ (solid line), $\alpha = \frac{\pi}{4}$ (dashed-dot), and $\alpha = \frac{\pi}{2}$ (dashed).

V. CONCLUSION

In this paper, we have studied the production of a charged Higgs boson in association with a W boson at the LHC in the Type III 2HDM with spontaneous CP violation. We find that the cross sections of $W^{\mp}H^{\pm}$ production are large enough to consider opportunity to observe these processes, and also find no observable effects of CP violation in these processes. In this model, the charged Higgs boson mass can be as low as about 150 GeV due to the effective complex Yukawa couplings, which distinguishes from the Type II 2HDM with charged Higgs mass being constrained to be larger than 350 GeV via rare B-meson decays. The possibility of detecting charged Higgs has been studied extensively in Type II 2HDM or MSSM with leptonic decay $H^+ \to \tau^+\nu$ or hadronic decay $H^+ \to t\bar{b}$ [16–22,22,23,29,30]. As our main interest is the effect of the general Yukawa couplings with the spontaneous CP-violation, we have considered only at the parton level, and will leave the inclusion of parton showering, hadronization, full simulation of the detector, etc. for future investigation. As the new physics inputs have large uncertainties in the parameter space, we have not tried to include any higher order corrections in the production cross section, and only included the contributions from $b\bar{b}$ annihilation and gg fusion to the lowest order. Using up-to-date information on the input parameters and proton parton distribution functions, we have presented theoretical predictions for the $W^{\pm}H^{\mp}$ production cross section. It has been found that unless very large, top-Yukawa coupling the gg fusion has small contributions and the $b\bar{b}$ annihilation is the dominated mechanism for $W^{\mp}H^{\pm}$ production at LHC. Apart from the fully integrated cross section, we have also analyzed distributions in p_T and considered the effect of the mixing between light CP-even h_1 and CP-odd h_3. As a consequence, it has been shown that the mixing effect is generally small unless the mass gap of the neutral Higgs becomes large.

ACKNOWLEDGMENTS

The authors would like to thank Ci Zhuang for useful discussions. Y. Tang would like to thank Y. Q. Ma for helpful discussions. This work was supported in part by the National Science Foundation of China (NSFC) under Grants Nos. 10821504 and 10975170, and the key Project of Knowledge Innovation Program (PKIP) of the Chinese Academy of Sciences.

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