Exclusive $B \to VV$ decays and CP violation in the general two-Higgs-doublet model

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Using the general factorization approach, we present a detailed investigation for the branching ratios, CP asymmetries and longitudinal polarization fractions in all charmless hadronic $B \to VV$ decays (except for the pure annihilation processes) within the most general two-Higgs-doublet model with spontaneous CP violation. It is seen that such a new physics model only has very small contributions to the branching ratios and longitudinal polarization fractions. However, as the model has rich CP-violating sources, it can lead to significant effects on the CP asymmetries, especially on those of penguin-dominated decay modes, which provides good signals for probing new physics beyond the SM in the future B-physics experiments.

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I. INTRODUCTION

During recent years, tremendous progress in B physics has been made through the fruitful interplay between theory and experiment. The precise measurements of the B-meson decays can provide an insight into very high energy scales via the indirect loop effects of new physics beyond the standard model (SM), which makes the study of exclusive nonleptonic B-meson decays of great interest.

In the SM, the phenomenon of CP violation can be accommodated in an efficient way through a complex phase entering the quark-mixing matrix, which governs the strength of the charged-current interactions of the quarks. This Kobayashi-Maskawa (KM) [1] mechanism of CP violation is the subject of detailed investigation in these few decades. However, its origin remains unknown as it is put into the standard model through the complex Yukawa couplings. Moreover, the baryon asymmetry of the universe requires new sources of CP violation. Many possible extensions of the SM in the Higgs sector have been proposed [2], and it was suggested that CP symmetry may break down spontaneously [3]. A consistent and simple model, which provides a spontaneous CP violation mechanism, has been constructed completely in a general two-Higgs-doublet model (2HDM) [4,5] without imposing the *ad hoc* discrete symmetry, which is now commonly called the type III 2HDM. The type III 2HDM, which allows flavor-changing neutral currents (FCNCs) at tree level but suppressed by approximate $U(1)$ flavor symmetry, has attracted much more interest. It is known that FCNCs are suppressed in low-energy experiments, especially for the lighter two generation quarks. Thus, the type III 2HDM can be parametrized in a way to satisfy the current experimental constraints. On the other hand, constraints on the general 2HDM from the neutral mesons mixing ($K^0 - \bar{K}^0$, $D^0 - \bar{D}^0$, and $B^0 - \bar{B}^0$) [6,7] and from the radiative decays of bottom quark [8] have also been studied in detail.

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In recent years, there have been many works about the B-meson decays within the two-Higgs-doublet model. In Refs. [9,10], the authors have studied the $B \to PP, PV$ decays (with P and V denoting the pseudoscalar and vector mesons, respectively) within the type III 2HDM. Since, through the measurements of magnitudes and phases of various helicity amplitudes, the charmless hadronic $B \to VV$ decay modes can reveal more dynamics of exclusive B decays than $B \to PP$ and $B \to PV$ decays, in the present work we are going to make a detailed study for $B \to VV$ decays within the type III 2HDM by emphasizing the new physics contributions. It will be seen that this specific new physics has remarkable effects on CP asymmetries, especially on the parameter S_f for the penguin-dominated decay modes. On the other hand, the new physics is found to have very small contributions to the branching ratios and the transverse polarizations. Furthermore, the polarization anomaly observed in $B \to \rho K^*$ and $B \to \phi K^*$ modes cannot be improved in our current considered parameter spaces.

The paper is organized as follows: In Sec. II, we first describe the theoretical framework, including a brief introduction for the two-Higgs-doublet model with spontaneous CP violation, the effective Hamiltonian, as well as the decay amplitudes and CP violation formulas, which are the basic tools to estimate the branching ratios and CP asymmetries of B-meson decays. In Sec. III, we list the Wilson coefficients and the other relevant input parameters. Our numerical predictions for the branching ratios, CP asymmetries, and longitudinal polarization fractions are presented in Sec. IV. Our conclusions are presented in the last section.

II. THEORETICAL FRAMEWORK

A. Outline of the two-Higgs-doublet model

Motivated solely from the origin of CP violation, a general two-Higgs-doublet model with spontaneous CP violation (type III 2HDM) has been shown to provide one of the simplest and most attractive models in understanding the origin and mechanism of *CP* violation at the weak scale. In such a model, there exists more physical neutral and charged Higgs bosons and rich *CP* violating sources from a single *CP* phase of the vacuum. These new sources of *CP* violation can lead to some significant phenomenological effects, which are promising to be tested by the future *B* factory and the LHCb experiments. In this paper, we shall focus on the phenomenological applications of the type III 2HDM on the two-body charmless hadronic $B \to VV$ decays.

The two complex Higgs doublets in the general 2HDM are generally expressed as [4,5,11,12]

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

The corresponding Yukawa Lagrangian is given as

$$\mathcal{L}_Y = \eta_{ija}\bar{\psi}_{i,L}\tilde{\Phi}_a U_{j,R} + \xi_{ija}\bar{\psi}_{i,L}\Phi_a D_{j,R} + \text{H.c.}, \tag{2}$$

where the parameters η_{ija} and ξ_{ija} are real, so that the Lagrangian is *CP* invariant. After the symmetry is spontaneously broken down

$$\langle\phi_1^0\rangle = v_1 e^{i\alpha_1}, \qquad \langle\phi_2^0\rangle = v_2 e^{\alpha_2}, \tag{3}$$

and the Goldstone particles have been eaten, the physical Higgs bosons are

$$H_1 = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v + \phi_1^0 \end{pmatrix}, \qquad H_2 = \frac{1}{\sqrt{2}}\begin{pmatrix} H^+ \\ \phi_2^0 + i\phi_3^0 \end{pmatrix}, \tag{4}$$

where $H^\pm$ are the charged scalar mass eigenstates, (ϕ_1^0, ϕ_2^0, ϕ_3^0) are generally not the mass eigenstates but can be expressed as linear combinations of the mass eigenstates (H, h, A).

Then the Yukawa part of the Lagrangian for physical particles can be written as

$$\mathcal{L}_Y = \eta_{ij}^U\bar{\psi}_{i,L}\tilde{H}_1 U_{j,R} + \eta_{ij}^D\bar{\psi}_{i,L}H_1 D_{j,R} + \xi_{ij}^U\bar{\psi}_{i,L}\tilde{H}_2 U_{j,R} + \xi_{ij}^D\bar{\psi}_{iL}H_2 D_{j,R} + \text{H.c.}, \tag{5}$$

where

$$\begin{aligned} \eta_{ij}^U &= \eta_{ij1}\cos\beta + \eta_{ij2}e^{-\delta}\sin\beta, \\ \xi_{ij}^U &= -\eta_{ij1}e^{-\delta}\sin\beta + \eta_{ij2}\cos\beta, \\ \eta_{ij}^D &= \xi_{ij1}\cos\beta + \xi_{ij2}e^{-\delta}\sin\beta, \\ \xi_{ij}^D &= -\xi_{ij1}e^{-\delta}\sin\beta + \xi_{ij2}\cos\beta, \end{aligned} \tag{6}$$

and these couplings η^U, η^D, ξ^U, ξ^D are generally complex, which means *CP* violation. According to the CKM mechanism, after diagonalizing the fermion terms' couplings η^U and η^D, the other couplings become

$$\begin{aligned} \mathcal{L}_Y = {}& \bar{U}_i\frac{m^U}{v}U_R(v + \phi_1^0) + \bar{D}_L\frac{m^D}{v}D_R(v + \phi_1^0) \\ &+ \bar{U}_L\tilde{\xi}^U U_R(\phi_2^0 + i\phi_3^0) + \bar{D}_L\hat{\xi}^U U_R H^- \\ &+ \bar{U}_L\hat{\xi}^D D_R H^+ + \bar{D}_L\tilde{\xi}^D D_R(\phi_2^0 + i\phi_3^0) + \text{H.c}, \end{aligned} \tag{7}$$

with

$$\begin{aligned} \tilde{\xi}^{U,D} &= (V_L^{U,D})^{-1}\xi^{U,D}V_R^{U,D}, \\ \hat{\xi}^U &= \tilde{\xi}^U V_{\text{CKM}}, \qquad \hat{\xi}^D = V_{\text{CKM}}\tilde{\xi}^D. \end{aligned} \tag{8}$$

The Yukawa couplings may be parametrized as follows:

$$\tilde{\xi}_{ij} = \lambda_{ij}\frac{\sqrt{m_i m_j}}{v}. \tag{9}$$

with v the vacuum expectation value $v = 246$ GeV.

B. Effective Hamiltonian and decay amplitudes of $B \to VV$ decays

Using the operator product expansion and the renormalization group equation, the low-energy effective Hamiltonian for charmless hadronic *B*-meson decays with $\Delta B = 1$ can be written as

$$\mathcal{H}_{\text{eff}} = \frac{G_F}{\sqrt{2}}\sum_{p=u,c}V_{pb}V_{pq}^*\Big(C_1Q_1^p + C_2Q_2^p + \sum_{i=3,\ldots,16}[C_iQ_i + C_i'Q_i']\Big) + \text{H.c.}, \tag{10}$$

where $C_i(\mu)(i = 1, \ldots, 16)$ are the Wilson coefficients that can be calculated by perturbative theory, and Q_i are the quark and gluon effective operators, with Q_{1-10} and Q_{11-16} coming from the SM and from the type III 2HDM, respectively. Their explicit forms are defined as follows (taking $b \to sq\bar{q}$ transition as an example) [13]

$$Q_1 = (\bar{s}u)_{V-A}(\bar{u}b)_{V-A}, \qquad Q_2 = (\bar{s}_i u_j)_{V-A}(\bar{u}_j b_i)_{V-A}, \qquad Q_{3(5)} = (\bar{s}b)_{V-A}\sum_q (\bar{q}q)_{V-(+)A},$$

$$Q_{4(6)} = (\bar{s}_i b_j)_{V-A}\sum_q (\bar{q}_j q_i)_{V-(+)A}, \qquad Q_{7(9)} = \frac{3}{2}(\bar{s}b)_{V-A}\sum_q e_q(\bar{q}q)_{V+(-)A}, \qquad Q_{8(10)} = \frac{3}{2}(\bar{s}_i b_j)_{V-A}\sum_q e_q(\bar{q}_j q_i)_{V+(-)A},$$

$$Q_{11(13)} = (\bar{s}b)_{S+P}\sum_q \frac{m_q\lambda^*_{qq}(\lambda_{qq})}{m_b}(\bar{q}q)_{S-(+)P}, \qquad Q_{12(14)} = (\bar{s}_i b_j)_{S+P}\sum_q \frac{m_q\lambda^*_{qq}(\lambda_{qq})}{m_b}(\bar{q}_j q_i)_{S-(+)P},$$

$$Q_{15} = \bar{s}\sigma^{\mu\nu}(1+\gamma_5)b\sum_q \frac{m_q\lambda_{qq}}{m_b}\bar{q}\sigma_{\mu\nu}(1+\gamma_5)q, \qquad Q_{16} = \bar{s}_i\sigma^{\mu\nu}(1+\gamma_5)b_j\sum_q \frac{m_q\lambda_{qq}}{m_b}\bar{q}_j\sigma_{\mu\nu}(1+\gamma_5)q_i, \tag{11}$$

where $(\bar{q}_1 q_2)_{V\pm A} = \bar{q}_1\gamma_\mu(1\pm\gamma_5)q_2$ and $(\bar{q}_1 q_2)_{S\pm P} = \bar{q}_1(1\pm\gamma_5)q_2$, with qu, d, s, c, b, and e_q is the electric charge number of q quark. The operators Q_i' in Eq. (11) are obtained from Q_i via exchanging $L \leftrightarrow R$, and we shall neglect their effects in our calculations for they are suppressed by a factor m_s/m_b in model III 2HDM. The Wilson coefficients $C_i(i = 1, \ldots, 10)$ have been calculated at leading order (LO) [14,15] and at next-to-leading order (NLO) [16] in the SM and also at LO in 2HDM [11], while $C_i(i = 11, \ldots, 16)$ at LO can be found in Refs. [13,17].

Having defined the effective Hamiltonian $H_{\rm eff}$ in terms of the four-quark operators Q_i, we can then proceed to calculate the hadronic matrix elements with the generalized factorization assumption [18–21] based on the naive factorization approach. It is known that the generalized factorization approach gives a solution to the scale problem in naive factorization. In this note, we extract the μ dependence from the matrix element $\langle O(\mu)\rangle$ and combine it with the μ-dependent Wilson coefficients function to form μ-dependent effective coefficients. Schematically, we may write

$$\langle \mathcal{H}_{\rm eff}\rangle = C(\mu)\langle O(\mu)\rangle = C(\mu)g(\mu)\langle O\rangle_{\rm tree} = C^{\rm eff}\langle O\rangle_{\rm tree}. \tag{12}$$

In principle, the Wilson coefficients $c^{\rm eff}$ should be renormalization scale independent. Thus it is necessary to incorporate QCD and EW corrections to the operators:

$$\langle O_i(\mu)\rangle = \left[\hat{\mathbf{1}} + \frac{\alpha_s(\mu)}{4\pi}\hat{\mathbf{m}}_s(\mu) + \frac{\alpha_{\rm em}}{4\pi}\hat{\mathbf{m}}_{\rm em}(\mu)\right]_{ij}\langle O\rangle_{\rm tree}, \tag{13}$$

with

$$c_i^{\rm eff}(\mu) = \left[\hat{\mathbf{1}} + \frac{\alpha_s(\mu)}{4\pi}\hat{\mathbf{m}}_s(\mu) + \frac{\alpha_{\rm em}}{4\pi}\hat{\mathbf{m}}_{\rm em}(\mu)\right]_{ij}c_j(\mu). \tag{14}$$

For simplicity, in this paper we will write $c_i^{\rm eff}$ as c_i in following. $\hat{\mathbf{m}}_s(\mu)$ and $\hat{\mathbf{m}}_{\rm em}(\mu)$ can be estimated through the absorptive parts of penguin topologies with up and charm quarks running in the loops. By doing this, we can get the CP-conserving strong phases, which are crucial for direct CP asymmetries.

For two-body charmless hadronic $B \to VV$ decays, the decay amplitude of the local four fermion operators is defined as

$$A_h \equiv \frac{G_F}{\sqrt{2}}\langle V_1(h_1)V_2(h_2)|(\bar{q}_2 q_3)_{V\pm A}(\bar{b}q_1)_{V-A}|B\rangle, \tag{15}$$

where h_1 and h_2 are the helicities of the final-state vector mesons V_1 and V_2 with four-momentum p_1 and p_2, respectively. Since the B meson has spin zero, in the rest frame of the B-meson system, the two vector mesons have the same helicity due to helicity conservation. Therefore three polarization states are possible in $B \to VV$ decays with one longitudinal (L) and two transverse, corresponding to helicities $h = 0$ and $h = \pm$ (here $h_1 = h_2 = h$), respectively. We define the three helicity amplitudes as follows

$$\begin{aligned} A_0 &= A(B \to V_1(p_1, \epsilon_1^0)V_2(p_2, \epsilon_2^0)), \\ A_\pm &= A(B \to V_1(p_1, \epsilon_1^\pm)V_2(p_2, \epsilon_2^\pm)). \end{aligned} \tag{16}$$

We choose the momentum $\vec{p}_2$ to be directed in the positive z-direction in the B-meson rest frame, and the polarization four-vectors of the light vector mesons such that in a frame where both light mesons have large momentum along the z-axis. They are given by

$$\epsilon_1^{\pm\mu} = \epsilon_2^{\mp\mu} = (0, \pm 1, i, 0)/\sqrt{2}, \qquad \epsilon_{1,2}^{0\mu} = p_{1,2}^\mu/m_{1,2}, \tag{17}$$

where m_1 and m_2 are the masses of V_1 and V_2 mesons, respectively. Using the definitions for decay constants and form factors [22], the tree-level hadronic matrix elements of the effective operators Q_i can be decomposed as the following two amplitudes

$$A_h = \mathcal{V}_h + \mathcal{T}_h, \tag{18}$$

with

$$\begin{aligned} \mathcal{V}_h &\equiv \langle V_1(p_1, \epsilon_1^h)|V - A|B\rangle\langle V_2(p_2, \epsilon_2^h)|V - A|0\rangle, \\ \mathcal{T}_h &\equiv \langle V_1(p_1, \epsilon_1^h)|\sigma^{\mu\nu}(1+\gamma^5)|B\rangle \\ &\quad \times \langle V_2(p_2, \epsilon_2^h)|\sigma_{\mu\nu}(1+\gamma^5)|0\rangle. \end{aligned} \tag{19}$$

Here, for simplicity, we have omitted the quark spinors in the corresponding current operators in the above definitions. The three polarization amplitudes for $\mathcal{V}_h$ and $\mathcal{T}_h$ can be further written as

$$\begin{aligned} \mathcal{V}_0 &= i f_{V_2}(m_B^2 - m_1^2 - m_2^2)A_0^{V_1}, \\ \mathcal{V}_\pm &= i f_{V_2} m_2\left[A_1^{V_1}(m_1 + m_B) \mp V^{V_1}\frac{2m_B|p_c|}{m_B + m_1}\right], \\ \mathcal{T}_0 &= 0, \\ \mathcal{T}_\pm &= 2i f_{V_2}^\perp[2T_1^{V_1} m_B|p_c| \mp T_2^{V_1}(m_B^2 - m_1^2)]. \end{aligned} \quad (20)$$

From the amplitude given by Eq. (18), the branching ratio for $B \to VV$ decays then reads

$$\mathrm{Br}(B \to VV) = \frac{\tau_B|p_c|}{8\pi m_B^2}(|A_0|^2 + |A_+|^2 + |A_-|^2), \quad (21)$$

where τ_B is the lifetime of the B meson, and p_c is the center of mass momentum of either final-state meson with

$$|p_c| = \frac{\sqrt{[m_B^2 - (m_1 + m_2)^2][m_B^2 - (m_1 - m_2)^2]}}{2m_B}. \quad (22)$$

In order to compare the relative size of the three different helicity amplitudes, we can define the longitudinal polarization fraction as

$$f_L = \frac{|A_0|^2}{|A_0|^2 + |A_+|^2 + |A_-|^2}, \quad (23)$$

which measures the relative strength of the longitudinal polarization amplitude in a given decay mode.

C. *CP*-violating asymmetries in $B \to VV$ decays

Since there are abundant *CP* violation sources in the two-Higgs-doublet model, it is also necessary and interesting for us to discuss *CP* asymmetries in $B \to VV$ decays.

First, for charged $B^\pm$-meson decays, there is only one simple type of *CP* violating asymmetry, which detects direct *CP* violation

$$\mathcal{A}_{CP} \equiv \frac{\Gamma(B^+ \to f^+) - \Gamma(B^- \to f^-)}{\Gamma(B^+ \to f^+) + \Gamma(B^- \to f^-)}. \quad (24)$$

For neutral B-meson decays, there is another type of *CP* violation coming from the mixing between $B_q^0 - \bar{B}_q^0$ (here $q = d$ or s)

$$\begin{aligned} |B_q^0(t)\rangle &= g_+(t)|B_q^0\rangle + \frac{q}{p}g_-(t)|\bar{B}_q^0\rangle, \\ |\bar{B}_q^0(t)\rangle &= \frac{p}{q}g_-(t)|B_q^0\rangle + g_+|\bar{B}_q^0\rangle. \end{aligned} \quad (25)$$

In this case, there are in general four amplitudes which can be expressed as [23–25]

$$\begin{aligned} A_f &= \langle f|H_{\rm eff}|B_q^0\rangle, & \bar{A}_f &= \langle f|H_{\rm eff}|\bar{B}_q^0\rangle, \\ \bar{A}_{\bar{f}} &= \langle \bar{f}|H_{\rm eff}|\bar{B}_q^0\rangle, & A_{\bar{f}} &= \langle \bar{f}|H_{\rm eff}|B_q^0\rangle. \end{aligned} \quad (26)$$

For the $B_d - \bar{B}_d$ and $B_s - \bar{B}_s$ systems, the following approximations can be made

$$\begin{aligned} &\text{both } B_d \text{ and } B_s \text{ systems: } \left|\frac{q}{p}\right| \sim 1; \\ &\text{only } B_d \text{ system: } \Delta\Gamma \sim 0. \end{aligned} \quad (27)$$

Using the decay amplitudes and the approximations listed in Eqs. (26) and (27), the time-dependent decay probabilities for the B_d system can then be written as

$$\begin{aligned} \Gamma(B_d^0(t) \to f) &= \frac{|A_f|^2(1 + |\lambda_f|^2)}{2} \\ &\quad \times e^{-\Gamma t}\{1 + C_f\cos(\Delta m t) - S_f\sin(\Delta m t)\}, \\ \Gamma(\bar{B}_d^0(t) \to f) &= \frac{|A_f|^2(1 + |\lambda_f|^2)}{2} \\ &\quad \times e^{-\Gamma t}\{1 - C_f\cos(\Delta m t) + S_f\sin(\Delta m t)\}, \end{aligned} \quad (28)$$

while for the B_s system, we have

$$\begin{aligned} \Gamma(B_s^0(t) \to f) &= \frac{|A_f|^2(1 + |\lambda_f|^2)}{2}e^{-\Gamma t}\left[\cosh\left(\frac{\Delta\Gamma t}{2}\right)\right. \\ &\quad + D_f\sinh\left(\frac{\Delta\Gamma t}{2}\right) + C_f\cos(\Delta m t) \\ &\quad \left. - S_f\sin(\Delta m t)\right], \\ \Gamma(\bar{B}_s^0(t) \to f) &= \frac{|A_f|^2(1 + |\lambda_f|^2)}{2}e^{-\Gamma t}\left[\cosh\left(\frac{\Delta\Gamma t}{2}\right)\right. \\ &\quad + D_f\sinh\left(\frac{\Delta\Gamma t}{2}\right) - C_f\cos(\Delta m t) \\ &\quad \left. + S_f\sin(\Delta m t)\right], \end{aligned} \quad (29)$$

where Γ is the average decay width, $\Delta\Gamma$ and Δm are the width and mass difference, respectively. The other quantities are defined as

$$\begin{aligned} \lambda_f &\equiv \frac{q}{p}\frac{\bar{A}_f}{A_f}, & D_f &\equiv \frac{2\mathrm{Re}(\lambda_f)}{1 + |\lambda_f|^2}, \\ C_f &\equiv \frac{1 - |\lambda_f|^2}{1 + |\lambda_f|^2}, & S_f &\equiv \frac{2\mathrm{Im}(\lambda_f)}{1 + |\lambda_f|^2}. \end{aligned} \quad (30)$$

From Eqs. (28) and (29), we can get:

$$\begin{aligned} \mathcal{A}_{CP}(B_d \to f) &= -C_f\cos\Delta m t + S_f\sin\Delta m t, \\ \mathcal{A}_{CP}(B_s \to f) &= \frac{-C_f\cos\Delta m t + S_f\sin\Delta m t}{\cosh(\frac{\Delta\Gamma t}{2}) + D_f\sinh(\frac{\Delta\Gamma t}{2})}. \end{aligned} \quad (31)$$

III. INPUT PARAMETERS

The theoretical predictions in our calculations depend on many input parameters, such as the Wilson coefficients, the CKM matrix elements, the hadronic parameters, and so on. Here we present all the relevant input parameters as follows.

TABLE I. The Wilson coefficients $\tilde{C}_{11,\ldots,16} = \frac{m_s \lambda_{ss}^{(*)}}{m_b} C_{11,\ldots,16}$ in $b \to s$ transition at $\mu = m_b = 4.2$ GeV in 2HDM.

	Case A	Case B	Case C
$\tilde{C}_{11}$	$-0.0085 + 0.012i$	$-0.0085 + 0.018i$	$-0.010 + 0.012i$
$\tilde{C}_{12}$	0	0	0
$\tilde{C}_{13}$	$-0.0030 - 0.0049i$	$-0.0052 - 0.0069i$	$-0.0029 - 0.0052i$
$\tilde{C}_{14}$	$-0.000\,060 - 0.000\,10i$	$-0.000\,11 - 0.000\,14i$	$-0.000\,059 - 0.000\,10i$
$\tilde{C}_{15}$	$0.000\,033 + 0.000\,055i$	$0.000\,058 + 0.000\,078i$	$0.000\,032 + 0.000\,059i$
$\tilde{C}_{16}$	$-0.000\,10 - 0.000\,17i$	$-0.000\,18 - 0.000\,24i$	$-0.0001 - 0.000\,18i$

It has been shown from $B^0_{d,s} - \bar{B}^0_{d,s}$ mixings that the parameters $|\lambda_{cc}|$ and $|\lambda_{ss}|$ in Eq. (11) can reach to be around 100 [26], while their phases are not well constrained. In our present work we simply fix the phases to be $\pi/4$, and this choice will not cause any trouble in our numerical results. For the parameters λ_{tt} and λ_{bb}, the constraints come mainly from the experiments for $B - \bar{B}$ mixing, $\Gamma(b \to s\gamma)$, $\Gamma(b \to c\tau\bar{\nu}_\tau)$, ρ_0, R_b, $B \to PV$, and the electric dipole moments (EDMS) of the electron and neutron [10,11,13,17,27]. Based on the above analyses, we choose the following three typical parameter spaces which are allowed by the present experiments and have been adopted for the $B \to PV$ decays [10]

$$\text{Case A}: |\lambda_{tt}| = 0.15; \qquad |\lambda_{bb}| = 50,$$
$$\text{Case B}: |\lambda_{tt}| = 0.3; \qquad |\lambda_{bb}| = 30,$$
$$\text{Case C}: |\lambda_{tt}| = 0.03; \qquad |\lambda_{bb}| = 100,$$

and $\theta_{tt} + \theta_{bb} = \pi/2$. For the Higgs masses and the Wilson coefficients of $C_{1,\ldots,10}$ corresponding to the SM, we use the results listed in the paper [10], while for the Wilson coefficients in the type III 2HDM, we redefine them as $\tilde{C}_{11,\ldots,16} = \frac{m_s \lambda_{ss}^{(*)}}{m_b} C_{11,\ldots,16}$ in order to compare the contributions from those operators in SM, here the factor $\frac{m_s \lambda_{ss}^{(*)}}{m_b}$ is associated with the operators in 2HDM, the numerical values for $\tilde{C}_{11,\ldots,16}$ are listed in Table I.

As for the CKM matrix elements, we shall use the Wolfenstein parametrization [28] with the values [25]: $A = 0.8533 \pm 0.0512$, $\lambda = 0.2200 \pm 0.0026$, $\bar{\rho} = 0.20 \pm 0.09$, and $\bar{\eta} = 0.33 \pm 0.05$.

For the hadronic parameters, the decay constants, and the form factors, we list them in Tables II and III, respectively.

IV. NUMERICAL RESULTS AND DISCUSSIONS

In this section, we shall classify the 28 channels of B^+, B^0 and B_s decays into two light vector mesons according to

TABLE II. The hadronic input parameters [25] and the decay constants taken from the QCD sum rules [29] and lattice theory [30].

$\tau_{B^\pm}$	τ_{B_d}	τ_{B_s}	M_{B_d}	M_{B_s}	m_b
1.638 ps	1.528 ps	1.472 ps	5.28 GeV	5.37 GeV	4.2 GeV
m_t	m_u	m_d	m_c	m_s	m_{ρ^0}
174 GeV	3.2 MeV	6.4 MeV	1.1 GeV	0.105 GeV	0.77 GeV
$m_{\rho^\pm}$	m_ω	m_ϕ	$m_{K^{*\pm}}$	$m_{K^{*0}}$	$\Lambda_{\rm QCD}$
0.77 GeV	0.782 GeV	1.02 GeV	0.892 GeV	0.896 GeV	225 MeV
f_ρ	f_ω	f_{K^*}	f_ϕ	f_ρ^T	f_ω^T
0.205 GeV	0.195 GeV	0.217 GeV	0.231 GeV	0.147 GeV	0.133 GeV
$f_{K^*}^T$	f_ϕ^T				
0.156 GeV	0.183 GeV				

TABLE III. The relevant $B \to V$ transition form factors at $q^2 = 0$ taken from the light-cone sum rules (LCSR) [31,32].

decay channel	V	A_0	A_1	A_2	T_1	T_3
$B \to \rho$	0.323	0.303	0.242	0.221	0.267	0.176
$B \to \omega$	0.293	0.281	0.219	0.198	0.242	0.155
$B \to K^*$	0.411	0.374	0.292	0.259	0.333	0.202
$B_s \to \bar{K}^*$	0.311	0.363	0.233	0.181	0.26	0.136
$B_s \to \phi$	0.434	0.474	0.311	0.234	0.349	0.175

the reliability of the calculation for various observables, which is motivated by the dominated contributing operators. We shall give our predictions for the branching ratios, the CP asymmetries, and the longitudinal polarization fractions both in the SM and in the 2HDM. Comparisons with the current experiment data, if possible, are also made.

Before moving to the detailed discussions, some general observations of new physics effects on $B \to VV$ decays should be made. As can be seen from Eqs. (9) and (11), the contributions of new physics operators $O_{11,\ldots,16}$ are always proportional to the factor m_q/v. Thus, they are severely suppressed for the first generation quarks. In this case, for $B \to \rho K^*$, ωK^*, $\rho\rho$, $\omega\rho$, $\omega\omega$ and $B_s \to \rho K^*$, ωK^*, $K^* K^*$, $\rho\phi$, $\omega\phi$ decay channels, we can safely ignore the contributions from those new operators. Note that the new physics still has effects on the Wilson Coefficients C_{1-10}. On the other hand, for $B \to \phi K^*$, $\phi\rho$, $\phi\omega$ and $B_s \to \phi K^*$, $\phi\phi$ decay channels, since these are all induced by $b \to qs\bar{s}$ ($q = d, s$) transitions, we could not ignore the new operators' contributions any more in this case. In the general factorization approach, it is impossible to produce a vector meson via the scalar and/or pseudoscalar currents from the vacuum state, and hence the new operators Q_{11} and Q_{13} have no contributions to $B \to VV$ decays. Moreover, from the results listed in Table I, it can be seen that all the contributing new operators $Q_{12,14,15,16}$ have only very small (even zero) Wilson coefficients. It is therefore expected that the new physics will have very small effects on the branching ratios and transverse amplitudes (hence on the transverse polarization fractions) of $B \to VV$ decays.

A. *CP*-averaged branching ratios and direct *CP* violation

According to different decay modes, we shall give our predictions for the branching ratios and direct CP violations one by one.

(i) Color-allowed tree-dominated decays. Our predictions for the CP-averaged branching ratios and the direct CP asymmetries are presented in Table IV. From the numerical results, we can see that the branching ratios are all at 10^{-5} order, and the direct CP asymmetries are all very small since the penguin amplitude contributions are much smaller than the ones from the tree diagrams. Most predictions within the SM are consistent with the current experiment data, and the new physics has very small effects on these types of decays.

(ii) Color-suppressed tree-dominated decays. The numerical results are given in Table V, it is interesting to note that the branching ratios will generally become smaller after including the new physics contributions except for the $B \to \rho^0\rho^0$ mode. Furthermore, there are big direct CP violations in

TABLE IV. The CP-averaged branching ratios (in unit of 10^{-6}) (first line) and the direct CP violations (second line) for the color-allowed tree-dominant processes both in the SM and in the type III 2HDM. Cases A-C stand for the three different parameter spaces listed in Sec. III.

Decay modes	Case A	Case B	Case C	SM	Exp.
$B^+ \to \rho^+\rho^0$	14.59	14.59	14.59	15.53	18.2 ± 3.0
	-0.004	-0.004	-0.004	-0.002	-0.08 ± 0.13
$B^0 \to \rho^+\rho^-$	26.33	25.93	26.73	27.49	$24.2^{+3.1}_{-3.2}$
	-0.043	-0.043	-0.042	-0.035	
$B^+ \to \rho^+\omega$	12.66	12.47	12.85	13.97	$10.6^{+2.6}_{-2.3}$
	-0.042	-0.043	-0.042	-0.034	0.04 ± 0.18
$B_s \to \rho^+ K^{*-}$	36.88	36.32	37.44	38.50	
	-0.043	-0.043	-0.042	-0.035	

TABLE V. The same as Table IV but for color-suppressed tree-dominant processes.

Decay modes	Case A	Case B	Case C	SM	Exp.
$B^0 \to \rho^0\rho^0$	0.0814	0.0897	0.0754	0.065	0.86 ± 0.28
	0.176	0.218	0.119	0.153	
$B^+ \to \omega\omega$	0.112	0.110	0.115	0.160	<4.0
	-0.117	-0.088	-0.144	-0.207	
$B_s \to \rho^0\bar{K}^{*0}$	0.081	0.090	0.073	0.092	$<7.67 \times 10^{-4}$
	0.176	0.218	0.119	0.153	
$B_s \to \omega\bar{K}^{*0}$	0.183	0.180	0.187	0.262	
	-0.167	-0.088	-0.144	-0.207	
$B^+ \to \rho^0\omega$	0.024	0.024	0.024	0.076	<1.5
	-0.063	-0.063	-0.063	-0.035	

these decay processes except for the $B^+ \to \rho^0 \omega$ mode, and the new physics has more effects on the direct CP asymmetries than on the branching ratios through the Wilson coefficient functions, although there are no new operator contributions to the hadronic matrix elements in this type decays within our approximations. Compared to Case A and Case C, Case B has the biggest corrections to the CP asymmetries of the SM.

(iii) Penguin-dominated decays. We may divide such decays into two types: $\Delta S = 1$ and $\Delta D = 1$ decay modes. They correspond to the upper and the lower parts in Table VI, respectively. From the numerical results, we can see that all 11 $\Delta S = 1$ decay modes have branching ratios up to 10^{-6} or even to 10^{-5} order, since they involve the relatively large CKM matrix elements V_{ts}^*, while the $\Delta D = 1$ ones have much smaller branching ratios of order of 10^{-7} due to the smaller CKM matrix elements V_{td}^*. For $B \to \omega K^*$ and $B_s \to \phi\phi$ decay modes, our predictions for the branching ratios with including the new operator contributions have similar results as the ones within the SM, which, however, are not quite consistent with the current experimental data; the numerical results for $B \to \omega K^*$ modes are larger than the current experiment limit, and the prediction for $B_s \to \phi\phi$ is about 2 times larger than the present data. For the other decay modes, our predictions for the branching ratios are in general agreement with the data. As for the direct CP asymmetries, there are big CP violations in some decay modes, and the new physics can lead to remarkable effects. Our predictions are consistent with the data in all these decay modes.

(iv) Electroweak penguin or QCD flavor singlet dominated decays. As can be seen from Table VII, these types of decays are expected to have smaller branching ratios due to the large cancellations among the different Wilson coefficients. Although there are new operator contributions in $B \to \rho\phi$ and $\omega\phi$ decay modes, the predicted branching ratios are still small. The direct CP asymmetries for these decays are all small, and the new physics effects on these observables are not prominent. Because of the lack of accurate experimental data, we could not compare our predictions with the data yet.

TABLE VI. The same as Table IV but for the penguin-dominated decay modes. The upper and the lower parts correspond to $\Delta S = 1$ and $\Delta D = 1$ processes, respectively.

Decay modes	Case A	Case B	Case C	SM	Exp.
$B^+ \to \rho^+ K^{*0}$	7.169	7.409	7.027	7.287	9.2 ± 1.5
	0.084	0.117	0.049	0.018	-0.01 ± 0.16
$B^+ \to \rho^0 K^{*+}$	5.853	6.229	5.526	5.575	<6.1
	0.184	0.196	0.169	0.122	$0.20^{+0.32}_{-0.29}$
$B^0 \to \rho^0 K^{*0}$	6.396	6.513	6.324	6.245	5.6 ± 1.6
	0.054	0.073	0.033	0.018	0.09 ± 0.19
$B^0 \to \rho^- K^{*+}$	6.046	6.738	5.445	5.571	<12
	0.295	0.301	0.283	0.199	
$B^0 \to \omega K^{*0}$	3.412	3.513	3.351	3.498	<2.7
	0.078	0.107	0.048	0.024	
$B^+ \to \omega K^{*+}$	3.247	3.5697	2.965	3.123	<3.4
	0.265	0.274	0.251	0.176	
$B^0 \to \phi K^{*0}$	9.276	9.704	9.221	9.318	9.5 ± 0.8
	0.045	0.081	-0.002	0.020	-0.01 ± 0.06
$B^+ \to \phi K^{*+}$	9.867	10.32	9.775	9.979	10.0 ± 1.1
	0.039	0.074	-0.013	0.020	-0.01 ± 0.08
$B_s \to \phi\phi$	28.99	30.34	28.64	28.85	$14^{+8}_{-7} \times 10^{-6}$
	0.054	0.089	0.006	0.020	
$B_s \to \bar{K}^{*0} K^{*0}$	9.303	9.614	9.118	9.456	$<1.681 \times 10^{-3}$
	0.084	0.117	0.049	0.018	
$B_s \to K^{*+} K^{*-}$	8.404	9.366	7.569	7.744	
	0.295	0.302	0.283	0.199	
$B^0 \to \bar{K}^{*0} K^{*0}$	0.410	0.420	0.413	0.408	$0.49^{+0.17}_{-0.14}$
	-0.092	-0.061	-0.133	-0.145	
$B^+ \to K^{*+} K^{*0}$	0.439	0.450	0.443	0.437	<2.2
	-0.092	-0.061	-0.133	-0.145	
$B_s \to \phi \bar{K}^{*0}$	0.517	0.532	0.521	0.526	$<1.013 \times 10^{-3}$
	-0.094	-0.056	-0.145	-0.161	

TABLE VII. The same as Table IV but for the electroweak penguin or QCD flavor singlet dominated decays.

Decay modes	Case A	Case B	Case C	SM	Exp.
$B^+ \to \rho^+\phi$	0.0054	0.0054	0.0054	0.0043	<16
	−0.011	−0.011	−0.011	−0.014	
$B^0 \to \rho^0\phi$	0.0025	0.0025	0.0025	0.0020	<13
	−0.011	−0.011	−0.011	−0.014	
$B^0 \to \omega\phi$	0.0022	0.0022	0.0022	0.0017	<1.2
	−0.011	−0.011	−0.011	−0.014	
$B_s \to \rho^0\phi$	0.796	0.796	0.796	0.687	$<6.17 \times 10^{-4}$
	0.0048	0.0048	0.0048	0.0039	
$B_s \to \phi\omega$	0.038	0.038	0.038	0.045	
	0.020	0.020	0.020	0.018	

(v) The pure annihilation decays. Only six decays belong to this class, namely $B^0 \to K^{*+}K^{*-}$, $B^0 \to \phi\phi$, $B_s \to \rho^+\rho^-$, $B_s \to \rho^0\rho^0$, $B_s \to \rho^0\omega$, and $B_s \to \omega\omega$. Because of the lack of the information for the $V_1 \to V_2$ transition form factor at large momentum transfers, we shall not consider them in detail in this paper.

B. Time-dependent *CP* violating parameters C_f, S_f and D_f

Since there are abundant *CP* violating sources in type III 2HDM, it is expected that there are relatively large *CP* violations in 2HDM than in the SM. Using the relevant formulas given in Sec. II, we can predict the time-dependent *CP* asymmetries in neutral B_d and B_s decays, with the numerical results given in Tables VIII and IX, respectively.

From these two tables, it is seen that for $B^0 \to \rho^+\rho^-$, $\rho^0\phi$ and $\omega\phi$ decay modes, the new physics has hardly any effect on the parameters C_f and S_f, even though there are new operators contributions in $B^0 \to \rho^0\phi$ and $\omega\phi$ decay modes. On the other hand, the new physics has remarkable effects on the other decay modes, especially on $B^0 \to \omega\omega$ one (for this mode the new physics can even change the sign of the parameter S_f). Furthermore, different parameter spaces also have remarkable effects on these *CP* violation parameters.

For the B_s system, there are new operator contributions only in the $B_s \to \phi\phi$ mode. As is expected, the new physics has remarkable influence on the parameters C_f, S_f, and D_f. For the other four decay modes, although there are no new operator contributions, the new physics still has big effects on the parameter S_f, but small effects on C_f and D_f.

C. The polarization in $B \to \rho K^*$ and ϕK^* decays

Motivated by the polarization anomaly observed by the *BABAR* [33], Belle [34], and CDF [35] experiments, we shall study the polarization in $B \to VV$ decays, especially in $B \to \rho K^*$ and ϕK^* decays in this section.

TABLE VIII. The time-dependent *CP* asymmetry parameters C_f (first line) and S_f (second line) for B_d decays both in the SM and in the type III 2HDM. Cases A-C stand for the three different parameter spaces listed in Sec. III.

Decay modes	Case A	Case B	Case C	SM
$B^0 \to \rho^+\rho^-$	0.043	0.043	0.042	0.035
	−0.95	−0.95	−0.95	−0.95
$B^0 \to \rho^0\rho^0$	−0.18	−0.22	−0.12	−0.15
	0.97	0.92	0.99	0.89
$B^0 \to \omega\rho^0$	0.063	0.063	0.063	0.029
	−0.61	−0.61	−0.62	−0.97
$B^0 \to \phi\rho^0$	0.011	0.011	0.011	0.014
	0.70	0.70	0.70	0.70
$B^0 \to \omega\phi$	0.011	0.011	0.011	0.014
	0.70	0.70	0.70	0.70
$B^0 \to \omega\omega$	0.12	0.09	0.14	0.21
	0.53	0.65	0.40	−0.18
$B^0 \to K^{*0}\bar{K}^{*0}$	0.092	0.061	0.13	0.15
	0.85	0.92	0.75	0.57

TABLE IX. The time-dependent CP asymmetry parameters C_f (first line), S_f (second line), and D_f (third line) for B_s decays both in the SM and in the type III 2HDM

Decay modes	Case A	Case B	Case C	SM
$B_s \to \phi\rho^0$	−0.005	−0.005	−0.005	−0.004
	0.052	0.052	0.052	0.14
	0.99	0.99	0.99	0.99
$B_s \to \phi\omega$	−0.020	−0.020	−0.020	−0.018
	0.23	0.23	0.23	0.49
	0.97	0.97	0.97	0.87
$B_s \to \phi\phi$	−0.054	−0.090	−0.060	−0.020
	0.33	0.49	0.14	−0.004
	0.94	0.87	0.99	1.0
$B_s \to K^{*+}K^{*-}$	−0.30	−0.30	−0.28	−0.20
	0.92	0.95	0.88	0.79
	0.25	0.12	0.39	0.57
$B_s \to \bar{K}^{*0}K^{*0}$	−0.085	−0.12	−0.049	−0.018
	0.31	0.45	0.15	−0.003
	0.95	0.88	0.99	1.0

One important point that should be noted is that the predictions for the branching ratios of $B \to \rho K^*$ and ϕK^* modes are well consistent with the experiment data, which means that if we want to solve the observed polarization anomaly, we need to find some way to reduce the longitudinal amplitude and enhance transverse ones simultaneously. Many studies have been made to try to provide possible resolutions to the anomaly both within the SM [36–39] and in various new physics models [40–42]. Here we only concentrate on the longitudinal polarization fraction and the main results are listed in Table X.

It is noted that the polarization anomaly could be well resolved by introducing the tensor operators $O_{T1} = \bar{s}\sigma^{\mu\nu}(1+\gamma^5)b\bar{s}\sigma_{\mu\nu}(1+\gamma_5)s$ and $O_{T8} = \bar{s}_i\sigma^{\mu\nu}(1+\gamma^5)b_j\bar{s}_j\sigma_{\mu\nu}(1+\gamma_5)s_i$ in Ref. [42]. It is interesting to see that these two operators have similar forms as Q_{15} and Q_{16} in Eq. (11). However, from the numerical results given by Table X, we can see that the predicted longitudinal polarization fraction f_L for these decay modes in the type III 2HDM is almost the same as the one within the SM. Although there are new operator contributions in $B \to \phi K^*$ modes, we still cannot resolve the polarization anomaly observed in this decay mode. This is due to the fact that the strength of new operators in 2HDM is severely suppressed by the factor $m_q\lambda_{qq}/m_b$. Moreover, as has already been mentioned in the beginning of this section, the Wilson coefficients of these new operators are very small, which also result in the small effects on the transverse amplitudes.

For the other $B \to VV$ decay modes, the predictions for longitudinal polarization fractions are always about 0.90–0.95. For simplicity, we shall not list the results in detail anymore.

In conclusion, adopting the current parameter spaces and taking the general factorization method which is obviously very crude, we could not resolve the polarization anomaly observed in $B \to \rho K^*$ and ϕK^* modes within the SM and 2HDM. The polarization anomalies may not be caused by new physics; it is more likely that one needs to have a better understanding for hadronic physics with an appropriate QCD approach.

V. CONCLUSIONS

Using the general factorization approach, we have studied all the $B \to VV$ decay modes except for pure annihilation decay channels both within the SM and in the two-Higgs-doublet model. From the numerical results given in the previous section, we can see that: for the branching ratios, our predictions are generally well consistent with the current experimental data expect for the $B_s \to \phi\phi$ decay mode, and the new physics has marginal or even negligible effects on this observable. However, the new physics can give remarkable contributions to the CP

TABLE X. The longitudinal polarization fractions f_L for $B \to \rho K^*$ and ϕK^* decay modes. Cases A-C stand for the three different parameter spaces in the type III 2HDM.

Decay modes	SM	Case A	Case B	Case C	Exp.
$B^+ \to \rho^+ K^{*0}$	0.91	0.91	0.91	0.91	0.48 ± 0.08
$B^0 \to \rho K^{*0}$	0.95	0.95	0.93	0.93	0.57 ± 0.12
$B^+ \to \phi K^*$	0.89	0.89	0.89	0.89	0.50 ± 0.05
$B^0 \to \phi K^{*0}$	0.89	0.89	0.89	0.89	0.491 ± 0.032

asymmetry parameters C_f and S_f, especially to S_f in the penguin-dominated decay modes. Unfortunately, our predictions for the longitudinal polarization fractions of $B \to \rho K^*$ and ϕK^* decay modes in 2HDM are still as large as the ones in the SM, which are much larger than the experimental data. Some new mechanisms may be needed to improve those discrepancies.

For simplicity, in this paper we have neglected the contributions from annihilation and exchange diagrams, although they may play a significant rule in some decay channels. In our numerical calculations, we have only considered three possible parameter spaces for the type III 2HDM. Also we have totally neglected the first generation Yukawa couplings and the off-diagonal matrix elements of the Yukawa coupling matrix, in order to eliminate the FCNC at tree level. However, it is possible that the FCNC involving the third generation quarks still exists at tree level, making the constraints less stronger. In a word, we do not exclude the possibility of improving the predictions by using the other factorization methods with the annihilation and exchange diagram contributions included, by choosing other parameters spaces, or even by introducing additional fourth-generation quarks [43].

In conclusion, we have shown that the new Higgs bosons in the type III 2HDM with spontaneous *CP* violation can have significant effects on some charmless $B \to VV$ decays, especially for the penguin-dominated decay modes, which can be used as good signals to test the SM and to explore new physics from more precise measurements in the future *B*-factory experiments.

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