Implications of family nonuniversal Z' model on $B \to K_0^*(1430)\pi$ decays

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Within the QCD factorization formalism, we study the possible impacts of the nonuniversal Z' model, which provides a flavor-changing neutral current at the tree level, on rare decays $B \to K_0^*(1430)\pi$. Under two different scenarios, Scenerio-1 (S1) and Scenerio-2 (S2), for identifying the scalar meson $K_0^*(1430)$, the branching ratios, CP asymmetries, and isospin asymmetries are calculated in both the standard model and the family nonuniversal Z' model. We find that the branching ratios and CP asymmetries are sensitive to weak annihilation. In the standard model, with $\rho_A = 1$ and $\phi_A \in [-30°, 30°]$, the branching ratios of S1 (S2) are smaller (larger) than the experimental data. Adding the contribution of the Z' boson in two different cases (Case I and Case II), for S1, the branching ratios are still far away from experiment. For S2, in Case II, the branching ratios become smaller and can accommodate the data; in Case I, although the center values are enhanced, they can also explain the data with large uncertainties. Similar conclusions are also reached for CP asymmetries. Our results indicate that S2 is more favored than S1, even after considering new physics effects. Moreover, if there exists a nonuniversal Z' boson, Case II is preferred. All results can be tested in the LHC-b experiment and forthcoming super-B factory.

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I. INTRODUCTION

Recently, with rich events in two B factories, measurements of B meson nonleptonic charmless decays involving scalar mesons have become available. Among these decays, the processes $B \to K_0^*(1430)\pi$ are attractive since they are dominantly induced by the flavor-changing neutral current (FCNC) transition $b \to s\bar{q}q (q = u, d, s)$. Such a transition forbidden at the tree level in the standard model (SM) is expected to be an excellent ground for testing SM and searching for new physics beyond SM. Therefore, many similar decay modes induced by FCNC have been explored widely in the literatures, such as $B \to K\pi$, $K\eta^{(\prime)}$, $\phi K^{(*)}$. The recent reviews can be found, for example, in Ref. [1]. For the concerned decay modes $B \to K_0^*(1430)\pi$, the latest world averaged branching ratios from Heavy Flavor Average Group [2] are listed as:

$$\begin{aligned} &BR(B^+ \to K_0^{*0}(1430)\pi^+) = (45.1 \pm 6.3) \times 10^{-6};\\ &BR(B^0 \to K_0^{*+}(1430)\pi^-) = (33.5^{+3.9}_{-3.8}) \times 10^{-6}; \qquad (1)\\ &BR(B^0 \to K_0^{*0}(1430)\pi^0) = (11.7^{+4.2}_{-3.8}) \times 10^{-6}. \end{aligned}$$

Direct CP asymmetries of above decays have also been measured recently by *BABAR* and Belle experiments, which will be shown in Sec. IV. As direct CP violation is sensitive to the strong phase involved in the decay process, the comparison between theory and experiment will offer us information on the strong phases necessary for producing the measured direct CP asymmetries. Comparing the predicted results of the SM [3] with experimental data, i.e. Eq. (1), we notice that the theoretical results cannot accommodate the data well even with large uncertainties. So, it is worth while to explore whether some new physics models could explain the data.

When discussing the B meson nonleptonic charmless decays, the hadronic matrix elements are required. In the past few years, several novel methods have been proposed to study matrix elements related to exclusive hadronic B decays, such as naive factorization [4], generalized factorization [5], the perturbative QCD (pQCD) method [6], QCD factorization (QCDF) [7], the soft collinear effective theory [8], and so on. Among these approaches, QCDF based on collinear factorization is a systematic framework to compute these matrix elements from QCD theory, and it holds in the heavy quark limit $m_b \to \infty$ and the heavy quark symmetry. Thus, we shall use QCDF approach in the following calculations.

Although the study of scalar meson spectrum has been an interesting topic for a long time, the underlying structure of the light scalar meson is still controversial until now. In the literature, there are many schemes for the classification of them. Here we present two typical scenarios to describe the scalar mesons [9]. Scenario-1 (S1) is the naive 2-quark model: the nonet mesons below 1 GeV are treated as the lowest lying states, and the ones near 1.5 GeV are the first orbitally excited states. In Scenario-2 (S2), the nonet mesons near 1.5 GeV are regarded as the lowest lying states, while the mesons below 1 GeV may be viewed as exotic states beyond the 2-quark model. Since the mass of $K_0^*(1430)$ is very near 1.5 GeV, thus it should be composed by two quarks in both S1 and S2, but the decay constants and distribution amplitudes are different in the different scenarios. Under above pictures, the two body nonleptonic B decays involving scalar mesons have been explored in both QCDF [3,10,11] and pQCD approaches [12–17].

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As stated before, $B \to K_0^*(1430)\pi$ decays are dominantly induced by FCNC $b \to s\bar{q}q$ transition, hence they are sensitive to new physics contributions even if they are suppressed by a large mass parameter which characterizes the new physics scale. To search for signals of new physics, a model-independent analysis is not suitable for the current status. It is the purpose of this work to show that a new physics effect of similar size can be obtained from some models with an extra Z' boson. Z' bosons are known to naturally exist in certain well-motivated extensions of the SM, such as the string theory [18], the grand unified theories [19], the little Higgs models [20], light U-boson model [21], by adding additional $\mathrm{U}(1)'$ gauge symmetry.

Most studies have assumed that the Z' gauge couplings are family universal [22], so that they remain diagonal even in the presence of fermion flavor mixing by the GIM mechanism. However, in some string models building [23], it is possible to have family nonuniversal Z' couplings, because of different constructions of the different families. Also, another way to induce the family nonuniversal Z' couplings is to introduce exotic fermions having $\mathrm{U}(1)'$ charges different from those of the SM fermions, as occurs in models with the E_6 grand unified group [24]. In this case, mixing of the right-handed ordinary and exotic quarks, all $\mathrm{SU}(2)_L$ singlets, gives rise to FCNCs mediated by a heavy Z'. Now, rightly or wrongly, these models provide a motivation to consider nonuniversal couplings. Thus, a motivated Z' model for low-energy systems is the so-called family nonuniversal Z' model [25], where the Z' couplings are affected by fermion mixing and are not diagonal in the mass basis. Nontrivial FCNC effects at the tree level mediated by the Z' therefore are induced, which play an important role in explaining the CP asymmetries in the current high energy experiments by introducing new weak phases. Much more extensive discussions of specific models and other implications, along with a more complete set of references, are given in several reviews [26–28].

In fact, the effects of Z' boson have been studied extensively in low-energy flavor physics phenomena, such as neutral meson (K, D, or B) mixing, B meson decays involving the $b \to s$ transition in particular [29–31], single top production [32], as well as the leptons decays [25,33]. Very recently, in Ref. [33], employing current experimental data and taking a model-independent approach, Chiang *et al.* performed a comprehensive study of constraints on both flavor-conserving and-violating leptonic Z' couplings and found the couplings are small. With those results, one can further constrain the couplings between Z' and quarks by studying the $B \to X_s l^+ l^-$, $B_s \to l^+ l^-$, and $B \to K^{(*)} l^+ l^-$ decays, which have been of great interest recently. Of course, these couplings could also been constrained by studying the nonleptonic decays such as $B \to K\eta^{(\prime)}$, $B \to K\pi$, $B \to \pi\pi$ and $B \to \phi K^{(*)}$, though many hadronic uncertainties are involved.

In this work, we will show the implications of the family nonuniversal Z' model on $B \to K_0^*(1430)\pi$ decays. The layout of this paper is as follows. In Sec. II, we firstly present the formulaes of $B \to K_0^*(1430)\pi$ in the SM within the QCDF approach, involving the effective Hamiltonian and the amplitudes. In Sec. III, we specify our flavor-changing Z' model, and how the effective Hamiltonian responsible for hadronic B decays is modified. The numerical results and discussions are given in Sect. IV. The conclusions are presented in the final section.

II. CALCULATION IN THE STANDARD MODEL

In the 2-quark picture of S1 and S2, the two kinds of decay constants of scalar meson S are defined by:

$$\langle S(p)|\bar{q}_2\gamma_\mu q_1|0\rangle = f_S p_\mu, \quad \langle S(p)|\bar{q}_2 q_1|0\rangle = m_S \bar{f}_S. \tag{2}$$

The vector decay constant f_S and the scale-dependent scalar decay constant $\bar{f}_S$ are related by equations of motion

$$\mu_S f_S = \bar{f}_S, \quad \text{with} \quad \mu_S = \frac{m_S}{m_2(\mu) - m_1(\mu)}, \tag{3}$$

where m_2 and m_1 are the running current quark masses. Therefore, contrary to the case of pseudoscalar one, the vector decay constant of the scalar meson, namely, f_S, will vanish in the SU(3) limit. In other words, the vector decay constant of $K_0^*(1430)$ is fairly small.

As for the scalar meson wave function, the twist-2 and twist-3 light-cone distribution amplitudes (LCDAs) for different components could be combined into a single matrix element:

$$\begin{aligned}\langle K_0^{*+}(p)|\bar{u}_\beta(z)s_\alpha(0)|0\rangle = \frac{1}{\sqrt{6}}\int_0^1 dx e^{ixp\cdot z}\Big\{&\not{p}\phi_{K_0^{*+}}(x) \\ &+ m_S\phi^S_{K_0^{*+}}(x) \\ &+ \frac{1}{6}m_S\sigma_{\mu\nu}p^\mu z^\nu \phi^\sigma_{K_0^{*+}}(x)\Big\}_{\alpha\beta}.\end{aligned} \tag{4}$$

The distribution amplitudes $\phi_{K_0^*}(x)$, $\phi^S_{K_0^*}(x)$, and $\phi^\sigma_{K_0^*}(x)$ are normalized as:

$$\begin{aligned}&\int_0^1 dx\phi_{K_0^*}(x) = \frac{f_{K_0^*}}{2\sqrt{6}}, \\ &\int_0^1 dx\phi^S_{K_0^*}(x) = \int_0^1 dx\phi^\sigma_{K_0^*}(x) = \frac{\bar{f}_{K_0^*}}{2\sqrt{6}},\end{aligned} \tag{5}$$

and $\phi^T_{K_0^*}(x) = \frac{1}{6}\frac{d}{dx}\phi^\sigma_{K_0^*}(x)$. The twist-2 LCDA can be expanded in the Gegenbauer polynomials:

$$\phi_S(x,\mu) = \frac{1}{\sqrt{2N_c}}\bar{f}_S(\mu)6x(1-x)\sum_{m=1}^\infty B_m(\mu)C_m^{3/2}(2x-1). \tag{6}$$

The decay constants and the Gegenbauer moments for twist-2 wave function in two different scenarios have been studied explicitly in Refs. [3,10] using the QCD sum rule approach. As for the explicit form of the Gegenbauer moments for the twist-3 wave functions, there exist few drawbacks in the theoretical calculation [34], thus we choice the asymptotic form for simplicity:

$$\phi_S^s = \frac{1}{\sqrt{2N_c}}\bar{f}_f, \qquad \phi_S^T = \frac{1}{\sqrt{2N_c}}\bar{f}_S(1-2x). \tag{7}$$

For the pion meson, the asymptotic forms for twist-2 and twist-3 distribution amplitudes are also adopted:

$$\begin{aligned} \phi_P(x) &= f_P 6x(1-x), \qquad \phi_P^p(x) = f_P, \\ \phi_P^\sigma(x) &= f_P 6x(1-x). \end{aligned} \tag{8}$$

The form factors of $B \to P, S$ transitions are defined by [4]:

$$\begin{aligned} \langle P(p')|V_\mu|B(p)\rangle &= \left(P_\mu - \frac{m_B^2 - m_P^2}{q^2}q_\mu\right)F_1^{BP}(q^2) \\ &\quad + \frac{m_B^2 - m_P^2}{q^2}q_\mu F_0^{BP}(q^2), \\ \langle S(p')|A_\mu|B(p)\rangle &= -i\left[\left(P_\mu - \frac{m_B^2 - m_S^2}{q^2}q_\mu\right)F_1^{BS}(q^2) \right. \\ &\quad \left. + \frac{m_B^2 - m_S^2}{q^2}q_\mu F_0^{BS}(q^2)\right], \end{aligned} \tag{9}$$

where $P_\mu = (p + p')_\mu$, $q_\mu = (p - p')_\mu$. Various form factors have been evaluated by utilizing the relativistic covariant light-front quark model [35], and the momentum dependence is fitted to a 3-parameter form

$$F(q^2) = \frac{F(0)}{1 - a(q^2/m_B^2) + b(q^2/m_B^2)^2}. \tag{10}$$

The parameters a and b relevant for our purposes are refereed to Ref. [35].

Although we concentrate on the study of new physics, the used notation for new interacting operators will be similar to those presented in the SM. Therefore, it is useful to introduce the effective operators of the SM. Thus, we describe the effective Hamiltonian for $b \to sq\bar{q}$ decays as

$$\begin{aligned} H_{\rm eff} &= \frac{G_F}{\sqrt{2}}\sum_{p=u,c}\lambda_p\Big[C_1(\mu)O_1^{(q)}(\mu) + C_2(\mu)O_2^{(q)}(\mu) \\ &\quad + \sum_{i=3}^{10}C_i(\mu)O_i(\mu)\Big], \end{aligned} \tag{11}$$

where $\lambda_q = V_{qb}V_{qs}^*$ are the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements and the operators O_1–O_{10} are defined as [36]

$$\begin{aligned} &O_1^{(q)} = (\bar{s}_\alpha q_\beta)_{V-A}(\bar{q}_\beta b_\alpha)_{V-A}, \qquad O_2^{(q)} = (\bar{s}_\alpha q_\alpha)_{V-A}(\bar{q}_\beta b_\beta)_{V-A}, \qquad O_3 = (\bar{s}_\alpha b_\alpha)_{V-A}\sum_q(\bar{q}_\beta q_\beta)_{V-A}, \\ &O_4 = (\bar{s}_\alpha b_\beta)_{V-A}\sum_q(\bar{q}_\beta q_\alpha)_{V-A}, \qquad O_5 = (\bar{s}_\alpha b_\alpha)_{V-A}\sum_q(\bar{q}_\beta q_\beta)_{V+A}, \qquad O_6 = (\bar{s}_\alpha b_\beta)_{V-A}\sum_q(\bar{q}_\beta q_\alpha)_{V+A}, \\ &O_7 = \frac{3}{2}(\bar{s}_\alpha b_\alpha)_{V-A}\sum_q e_q(\bar{q}_\beta q_\beta)_{V+A}, \qquad O_8 = \frac{3}{2}(\bar{s}_\alpha b_\beta)_{V-A}\sum_q e_q(\bar{q}_\beta q_\alpha)_{V+A}, \\ &O_9 = \frac{3}{2}(\bar{s}_\alpha b_\alpha)_{V-A}\sum_q e_q(\bar{q}_\beta q_\beta)_{V-A}, \qquad O_{10} = \frac{3}{2}(\bar{s}_\alpha b_\beta)_{V-A}\sum_q e_q(\bar{q}_\beta q_\alpha)_{V-A}, \end{aligned} \tag{12}$$

with α and β being the color indices. In Eq. (11), O_1–O_2 are from the tree level of weak interactions, O_3–O_6 are the so-called QCD penguin operators and O_7–O_{10} are the electroweak (EW) penguin operators, while C_1–C_{10} are the corresponding Wilson coefficients.

In the QCDF approach, the contribution of the nonperturbative sector is dominated by the form factors and the nonfactorizable impact in the hadronic matrix elements is controlled by hard gluon exchange. The hadronic matrix elements of the decay can be written as

$$\begin{aligned} \langle M_1M_2|O_i|B\rangle &= \sum_j F_j^{B\to M_1}\int_0^1 dx T_{ij}^I(x)\Phi_{M_1}(x) \\ &\quad + \int_0^1 d\xi\int_0^1 dx\int_0^1 dy T_i^{II}(\xi,x,y) \\ &\quad \times \Phi_B(\xi)\Phi_{M_1}(x)\Phi_{M_2}(y). \end{aligned} \tag{13}$$

Here T_{ij}^I and T_i^{II} denote the perturbative short-distance interactions and can be calculated perturbatively. $\Phi_X(x)$ are nonperturbative light-cone distribution amplitudes, which should be universal. Using the weak effective Hamiltonian given by Eq. (11) and the definitions of a_i and b_i in Refs. [3,7], we can now write the decay amplitudes of $B \to K_0^*(1430)\pi$ as:

$$\begin{aligned} &A(B^- \to \bar{K}_0^{*0}(1430)\pi^-) \\ &\quad = \frac{G_F}{\sqrt{2}}\sum_{p=u,c}\lambda_p\Big\{\Big(a_4^p - r_\chi^{K_0^*}a_6^p - \frac{1}{2}(a_{10}^p - r_\chi^{K_0^*}a_8^p)\Big)_{\pi K_0^*} \\ &\qquad \times f_{K_0^*}F_0^{B\pi}(m_{K_0^*}^2)(m_B^2 - m_\pi^2) \\ &\qquad + f_B(b_2\delta_u^p + b_3 + b_{3,\rm EW})_{\pi K_0^*}\Big\}, \end{aligned} \tag{14}$$

$$A(B^- \to K_0^{*-}(1430)\pi^0) = \frac{G_F}{2}\sum_{p=u,c}\lambda_p\Big\{(a_1\delta_u^p + a_4^p - r_\chi^{K_0^*}a_6^p + a_{10}^p - r_\chi^{K_0^*}a_8^p)_{\pi K_0^*} f_{K_0^*}F_0^{B\pi}(m_{K_0^*}^2)(m_B^2 - m_\pi^2) - \Big[a_2\delta_u^p + \frac{3}{2}(a_9 - a_7)\Big]_{K_0^*\pi} f_\pi F_0^{BK_0^*}(m_\pi^2)(m_B^2 - m_{K_0^*}^2) + f_B(b_2\delta_u^p + b_3 + b_{3,\rm EW})_{\pi K_0^*}\Big\}, \tag{15}$$

$$A(\bar B^0 \to K_0^{*-}\pi^+) = \frac{G_F}{\sqrt{2}}\sum_{p=u,c}\lambda_p\Big\{(a_1\delta_u^p + a_4^p - r_\chi^{K_0^*}a_6^p + a_{10}^p - r_\chi^{K_0^*}a_8^p)_{\pi K_0^*} f_{K_0^*}F_0^{B\pi}(m_{K_0^*}^2)(m_B^2 - m_\pi^2) + f_B\Big(b_3 - \frac{1}{2}b_{3,\rm EW}\Big)_{\pi K_0^*}\Big\}, \tag{16}$$

$$A(\bar B^0 \to \bar K_0^{*0}\pi^0) = \frac{G_F}{2}\sum_{p=u,c}\lambda_p\Big\{\Big(-a_4^p + r_\chi^{K_0^*}a_6^p + \frac{1}{2}(a_{10}^p - r_\chi^{K_0^*}a_8^p)\Big)_{\pi K_0^*} \times f_{K_0^*}F_0^{B\pi}(m_{K_0^*}^2)(m_B^2 - m_\pi^2) - \Big[a_2\delta_u^p + \frac{3}{2}(a_9 - a_7)\Big]_{K_0^*\pi} f_\pi F_0^{BK_0^*}(m_\pi^2)(m_B^2 - m_{K_0^*}^2) + f_B\Big(-b_3 + \frac{1}{2}b_{3,\rm EW}\Big)_{\pi K_0^*}\Big\}, \tag{17}$$

where $\lambda_p \equiv V_{pb}V_{ps}^*$ and

$$r_\chi^{K_0^*}(\mu) = \frac{2m_{K_0^*}^2}{m_b(\mu)(m_s(\mu) - m_q(\mu))}. \tag{18}$$

In the above formulas, the order of the arguments of the $a_i^p(M_1M_2)$ and $b_i(M_1M_2)$ coefficients is dictated by the subscript M_1M_2, where M_2 is the emitted meson and M_1 shares the same spectator quark with the B meson. For the annihilation diagram, M_1 is referred to the one containing an antiquark from the weak vertex, while M_2 contains a quark from the weak vertex. Note that the coefficients a_i come from vertex corrections and hard spectator corrections, and b_i represent of contribution of annihilation diagrams. Both a_i and b_i can be found in Ref. [3]. It must be emphasized that we shall evaluate the vertex corrections to the decay amplitudes at the scale $\mu = m_b/2$. In contrast, the hard spectator and annihilation contributions should be evaluated at the hard-collinear scale $\mu_h = \sqrt{\mu\Lambda_h}$ with $\Lambda_h \approx 500$ MeV.

In QCDF approach, the annihilation amplitude has endpoint divergences even at twist-2 level and the hard spectator scattering diagram at twist-3 order is power suppressed and posses soft and collinear divergences arising from the soft spectator quark. Since the treatment of endpoint divergences is model dependent, subleading power corrections generally can be studied only in a phenomenological way. We shall follow [3,7] to parameterize the endpoint divergence $X_A \equiv \int_0^1 dx/\bar x$ in the annihilation diagram as

$$X_A = \ln\Big(\frac{m_B}{\Lambda_h}\Big)(1 + \rho_A e^{i\phi_A}), \tag{19}$$

with the unknown real parameters ρ_A and ϕ_A. Likewise, the endpoint divergence X_H in the hard spectator contributions can be parameterized in a similar manner. In the Sec. IV, we will see that such divergence is the main source of the uncertainty for the concerned decay modes.

III. THE FAMILY NON-UNIVERSAL Z' MODEL

As mentioned before, a family nonuniversal Z' model leads to FCNC at the tree level due to the nondiagonal chiral coupling matrix, which makes itself become interesting in some penguin dominate processes. The basic formalism of flavor-changing effects in the Z' model with family nonuniversal and/or nondiagonal couplings has been laid out in Refs. [25,26], to which we refer readers for detail. The detailed phenomenological analysis for various low-energy physics, especially for B meson decays, could be found in Refs. [29–31]. Here we just briefly review the ingredients needed in this paper.

In practice, neglecting the renormalization group (RG) running between m_W and $m_{Z'}$ and mixing between Z' and Z boson of the SM, we write the Z' term of the neutral-current Lagrangian in the gauge basis as

$$\mathcal{L} = -g' J'_\mu Z'^\mu, \tag{20}$$

where g' is the gauge coupling constant of extra $U(1)'$ group at the electroweak m_W scale. The chiral current J'_μ is expressed as:

$$J'_\mu = \bar\psi_i\gamma_\mu[(B_{ij}^L)P_L + (B_{ij}^R)P_R]\psi_j, \tag{21}$$

where the chirality projection operators are $P_{L,R} \equiv (1 \pm \gamma_5)/2$ and B_{ij}^X refers to the effective Z' couplings to the quarks i and j at the electroweak scale. For simplicity, we assume that the right hand couplings are flavor-diagonal and neglect B_{sb}^R. Compared with Eq. (11), the effective Hamiltonian for $b \to s\bar q q$ transition with Z' boson can be written as

$$\mathcal{H}_{\rm eff}^{Z'} = \frac{2G_F}{\sqrt{2}}\Big(\frac{g'm_Z}{g_1 m_{Z'}}\Big)^2 B_{sb}^L(\bar s b)_{V-A}\sum_q[B_{qq}^L(\bar q q)_{V-A} + B_{qq}^R(\bar q q)_{V+A}] + \text{h.c.}, \tag{22}$$

where $m_{Z'}$ is the mass of the new gauge boson. In fact, the forms of 4-quark operators in Eq. (22) already exist in the SM, so we rewrite it as

$$\mathcal{H}_{\text{eff}}^{Z'} = -\frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^* \sum_q (\Delta C_3 O_3^q + \Delta C_5 O_5^q + \Delta C_7 O_7^q + \Delta C_9 O_9^q) + \text{h.c.}, \tag{23}$$

where $O_i^q (i = 3, 5, 7, 9)$ are the effective 4-quark operators in the SM. ΔC_i denote the modifications to the corresponding SM Wilson coefficients, which are expressed as

$$\begin{aligned} \Delta C_{3,5} &= -\frac{2}{3V_{tb}V_{ts}^*}\left(\frac{g' m_Z}{g_1 m_{Z'}}\right)^2 B_{sb}^L (B_{uu}^{L,R} + 2B_{dd}^{L,R}), \\ \Delta C_{9,7} &= -\frac{4}{3V_{tb}V_{ts}^*}\left(\frac{g' m_Z}{g_1 m_{Z'}}\right)^2 B_{sb}^L (B_{uu}^{L,R} - B_{dd}^{L,R}), \end{aligned} \tag{24}$$

Generally, the diagonal elements of the effective coupling matrices $B_{qq}^{L,R}$ are expected to be real as a consequence of the hermiticity of the effective weak Hamiltonian. However, the off-diagonal one B_{sb} perhaps contains a new weak phase ϕ_s. We also suppose $B_{qq}^L = B_{qq}^R = B_{qq}$, so as to reduce the new parameters. For convenience we can represent ΔC_i as

$$\Delta C_{3,5} = 2\frac{|V_{tb}V_{ts}^*|}{V_{tb}V_{ts}^*}\zeta e^{i\phi_s}, \qquad \Delta C_{9,7} = 4\frac{|V_{tb}V_{ts}^*|}{V_{tb}V_{ts}^*}\xi e^{i\phi_s}, \tag{25}$$

where ζ and ξ are defined, respectively, as

$$\begin{aligned} \zeta &= -\frac{1}{3}\left(\frac{g' M_Z}{g_1 M_{Z'}}\right)^2 \left|\frac{B_{sb}^L}{V_{tb}V_{ts}^*}\right| (B_{uu} + 2B_{dd}), \\ \xi &= -\frac{1}{3}\left(\frac{g' M_Z}{g_1 M_{Z'}}\right)^2 \left|\frac{B_{sb}^L}{V_{tb}V_{ts}^*}\right| (B_{uu} - B_{dd}). \end{aligned} \tag{26}$$

While in general we can have a Z' contribution to the QCD penguins $\Delta C_{3,5}$ as well as the EW penguins $\Delta C_{9,7}$. It is stressed that the other SM Wilson coefficients may also receive contributions from the Z' boson through RG evolution. With our assumption that no significant RG running effect between M_Z' and M_W scales, the RG evolution of the modified Wilson coefficients is exactly the same as the ones in the SM [36].

In order to show the effects of Z' boson clearly, our analysis are divided into the two cases with two different simplifications,

$$\begin{aligned} &B_{uu} = -2B_{dd}, \quad \zeta = 0, \quad \xi = X, \quad \text{Case–I}; \\ &B_{uu} = B_{dd}, \quad \zeta = -X, \quad \xi = 0, \quad \text{Case–II}, \end{aligned} \tag{27}$$

with

$$X = \left(\frac{g' M_Z}{g_1 M_{Z'}}\right)^2 \left|\frac{B_{sb}^L B_{dd}}{V_{tb}V_{ts}^*}\right| = y \left|\frac{B_{sb}^L B_{dd}}{V_{tb}V_{ts}^*}\right|. \tag{28}$$

For Case I, we assume $B_{uu} = -2B_{dd}$ so that new physics is primarily manifest in the EW penguins. (The same assumption has been used widely in Refs. [29,30].) The Case II means that the new physics effect is similar to the QCD penguins. In all cases, the relations between B_{uu} and B_{dd} can be realized by setting a small mixing angle between Z and Z'. Thus, there are only two parameters, X and weak phase ϕ_s left, in the sequential numerical calculations and discussions.

IV. NUMERICAL RESULTS AND DISCUSSIONS

To obtain the numerical results, we list the parameters related to the SM firstly. As stated in Sec. I, because we have not a clear conclusion whether $K_0^*(1430)$ belongs to the first orbitally excited state (S1) or the low lying state (S2), we have to calculate the processes under both scenarios. So, the decay constants, Gegenbauer moments, and form factors in different scenarios are listed as follows [3]:

$$\begin{aligned} &\textbf{S1}: \ \bar{f}_{K_0^*}(1.0\ \text{GeV}) = -300\ \text{MeV}; \\ &\bar{f}_{K_0^*}(2.1\ \text{GeV}) = -370\ \text{MeV}; \qquad B_1(1.0\ \text{GeV}) = 0.58; \\ &B_1(2.1\ \text{GeV}) = 0.39; \qquad B_3(1.0\ \text{GeV}) = -1.20; \\ &B_3(2.1\ \text{GeV}) = -0.70; \qquad F_0^{BK_0^*}(0) = F_1^{BK_0^*}(0) = 0.21; \end{aligned} \tag{29}$$

$$\begin{aligned} &\textbf{S2}: \ \bar{f}_{K_0^*(1430)}(1.0\ \text{GeV}) = 445\ \text{MeV}; \\ &\bar{f}_{K_0^*(1430)}(2.1\ \text{GeV}) = 550\ \text{MeV}; \\ &B_1(1.0\ \text{GeV}) = -0.57; \qquad B_1(2.1\ \text{GeV}) = -0.39; \\ &B_3(1.0\ \text{GeV}) = -0.42; \qquad B_3(2.1\ \text{GeV}) = -0.25; \\ &F_0^{BK_0^*}(0) = F_1^{BK_0^*}(0) = 0.26. \end{aligned} \tag{30}$$

Now that the uncertainties for the above parameters have been explored explicitly in Ref. [3], and we will not discuss the errors caused by them in the current work.

In Ref. [3], the authors concluded that the theoretical errors are dominated by the $1/m_b$ power corrections due to the weak annihilations. Moreover, the weak annihilation contributions to $B \to SP$ could be much larger than the $B \to PP$ case, because the helicity suppression appeared in the $B \to PP$ case can be alleviated in the scalar production with the nonvanishing orbital angular momentum in the scalar state. In order to accommodate the data, one has to take into account the power corrections due to the ρ_H and ρ_A from the hard spectator interactions and weak annihilations, respectively. In Ref. [3], Cheng *et al.* found that the predictions are far away from the experimental data if by setting $\rho_A = 0$, which indicates that ρ_A will be nonzero. Meanwhile, for $B \to PP$, PV modes [7], the errors due to weak annihilations are comparable to or much smaller than the center values, and the fitting results show that $\rho_A = 1$ and $\phi_A = 0°$. Hence, in this work, we adopt $\rho_H = \rho_A = 1$, and set the strong phases $\phi_{A,H}$ in the ranges $[-30°, 30°]$.

With above parameters, we present our predictions of the SM in Table. I under two different scenarios. For the center values, we also assign $\phi_A = \phi_H = 0$. In order to

TABLE I. Branching ratios (in units of 10^{-6}) of $B \to K_0^*(1430)\pi$ in the SM and the nonuniversal Z' model.

	S1			S2			
Decay Mode	SM	Case I	Case II	SM	Case I	Case II	Expt
$B^- \to \bar{K}_0^{*0}(1430)\pi^-$	$23.0^{+1.2}_{-5.9}$	$25.7^{+5.0+2.8}_{-4.7-7.7}$	$17.2^{+1.3+19.6}_{-5.1-5.0}$	$74.7^{+1.0}_{-20.6}$	$93.8^{+1.7+20.6}_{-25.8-50.6}$	$53.8^{+1.3+16.5}_{-17.5-17.4}$	$45.1^{+6.3}_{-6.3}$
$B^- \to K_0^{*-}(1430)\pi^0$	$9.3^{+1.0}_{-1.9}$	$17.9^{+0.8+11.4}_{-3.3-17.3}$	$6.8^{+1.1+8.7}_{-1.6-2.2}$	$38.9^{+0.4}_{-8.9}$	$75.3^{+0.0+46.8}_{-14.5-73.1}$	$28.2^{+0.5+8.4}_{-7.7-8.9}$	
$\bar{B}^0 \to K_0^{*-}(1430)\pi^+$	$21.3^{+0.7}_{-5.1}$	$27.2^{+0.2+6.8}_{-6.4-15.6}$	$15.9^{+0.8+18.4}_{-4.4-4.6}$	$70.0^{+0.6}_{-17.2}$	$83.3^{+0.0+13.4}_{-16.2-36.3}$	$50.2^{+0.9+15.6}_{-14.6-16.4}$	$33.5^{+3.9}_{-3.8}$
$\bar{B}^0 \to \bar{K}_0^{*0}(1430)\pi^0$	$12.9^{+0.3}_{-3.6}$	$9.4^{+0.3+12.0}_{-2.7-2.9}$	$9.8^{+0.3+10.3}_{-3.1-2.7}$	$33.6^{+0.4}_{-9.8}$	$22.0^{+0.5+41.9}_{-7.2-8.5}$	$23.8^{+0.5+7.7}_{-8.2-8.0}$	$11.7^{+4.2}_{-3.8}$

obtain the errors, we scan randomly the points in the ranges $\phi_A \in [-30°, 30°]$ and $\phi_H \in [-30°, 30°]$. So, the only theoretical errors of the SM results are due to the strong phases ϕ_A and ϕ_H. Because we fully consider the weak annihilations, our results are much larger than those in Ref. [3], especially for the center values. Compared with the data, the theoretical results in this work are still much smaller (larger) than the data under two scenarios, except for mode $\bar{B}^0 \to \bar{K}_0^{*0}(1430)\pi^0$. If one wants to fit the data absolutely, $\rho_A \approx 1.3$ for S1 and $\rho_A \approx 0.7$ for S2 are required, respectively, which are a bit larger/smaller by 30% than the fitted results from $B \to PP, PV$. Compared with predictions of Ref. [15] obtained in the pQCD approach based on k_T factorization, our results are a bit larger than theirs in S2, but agree with their results in S1 with large uncertainties.

We next turn to the implications of the nonuniversal Z' model for the $B \to K_0^*(1430)\pi$ decays. Let us firstly consider the range of X, which is the most important parameter in this model. Generally, we always expect $g'/g_1 \sim 1$, if both the $U(1)$ gauge groups have the same origin from some grand unified theories. $M_Z/M_{Z'} \sim 0.1$ for TeV scale neutral Z' boson is also expected so that the Z' could be detected in the running of the Large Hadron Collider (LHC), which results in $y \sim 10^{-2}$. In the first paper of Refs. [29], assuming a small mixing between $Z - Z'$ bosons, the value of y is taken as $y \sim 10^{-3}$. In order to explain the mass difference of $B_s - \bar{B}_s$ mixing, $|B_{sb}^L| \sim |V_{tb}V_{ts}^*|$ is required. The experimental data of $B \to \phi K^{(*)}$, $K\pi$, $\pi\pi$ requires $|B_{sb}^L B_{ss}^{L,R}| \sim |V_{tb}V_{ts}^*|$, which indicates $|B_{qq}^L| \sim 1$. Above issues have been discussed widely in Refs. [29,30]. Summing up the above analysis, we thereby assume that $X \in (10^{-3}, 10^{-2})$. For weak phase ϕ_s, though many attempts have been done to constrain it [31], we here left it as a free parameter.

The calculated results for branching ratios with two different cases in the family nonuniversal Z' model are also exhibited in Table. I, and for the center values we use $X = 0.005$ and $\phi_s = 0°$. To obtain the second errors, we also scan randomly the points in the ranges $X \in [0.001, 0.01]$ and $\phi_s \in [-180°, 180°]$, while the first errors come from the weak annihilations. The table shows to us that the two cases of Z' models can change the branching ratios remarkably in the two different scenarios. It is clear that the Z' will enhance the branching ratios in Case I, while in Case II the branching ratios are decreased. The reason is that the variation tendencies of Wilson coefficients are different in the two different cases, which could be seen in Eq. (25) and (27) easily.

For S1, the branching ratios of the first three decay modes cannot agree with data unless the upper limits in Case II of the Z' model are taken. Unfortunately, with the upper limit values, the branching ratio of $\bar{B}^0 \to \bar{K}_0^{*0}(1430)\pi^0$ is much larger than the experimental data. For S2, the branching ratios with a Z' boson can accommodate experimental data well in two cases with large uncertainties. If we care about the center values very much, it seems that results of Case II are preferable. If further theories and/or experiments can confirm the existence of Z', one could correspondingly cross-check the couplings and the mass of it with all above results in turn.

In the experimental side, another important observable in B physics is CP asymmetry, in particular, of the direct CP asymmetry. In Table. II, we list the direct CP asymmetries of concerned modes in different scenarios and different cases of the Z' model. Generally, the strong phases calculable in the QCD factorization are so small that the CP asymmetries are at most a few percent, as shown in the table. In S1 we note that the center values have different signs with the experimental data. Adding the Z' contribution, although the large uncertainties perhaps

TABLE II. *CP* asymmetry (in %) of $B \to K_0^*(1430)\pi$ in the SM and the nonuniversal Z' model.

	S1			S2			
Decay mode	SM	Case I	Case II	SM	Case I	Case II	Expt
$B^- \to \bar{K}_0^{*0}(1430)\pi^-$	$1.0^{+1.9}_{-1.9}$	$1.0^{+1.4+0.2}_{-1.6-0.1}$	$1.2^{+2.4+0.2}_{-2.1-0.4}$	$0.06^{+0.6}_{-0.7}$	$0.06^{+0.4+0.3}_{-0.5-0.3}$	$0.03^{+0.66+0.03}_{-0.82-0.07}$	-5^{+5}_{-8}
$B^- \to K_0^{*-}(1430)\pi^0$	$-0.5^{+3.8}_{-2.6}$	$-0.4^{+2.7+0.6}_{-2.0-2.8}$	$-0.6^{+4.6+0.4}_{-2.8-0.2}$	$1.0^{+2.4}_{-2.9}$	$0.7^{+1.8+5.4}_{-2.0-1.2}$	$1.2^{+2.7+0.3}_{-3.5-0.2}$	
$\bar{B}^0 \to K_0^{*-}(1430)\pi^+$	$2.0^{+2.7}_{-3.9}$	$1.7^{+2.3+1.5}_{-2.6-0.7}$	$2.4^{+3.1+1.0}_{-4.6-1.4}$	$-0.8^{+2.9}_{-2.7}$	$-0.7^{+2.1+0.6}_{-2.2-0.9}$	$-1.1^{+3.6+0.2}_{-3.3-0.5}$	-7^{+14}_{-14}
$\bar{B}^0 \to \bar{K}_0^{*0}(1430)\pi^0$	$3.1^{+2.6}_{-4.0}$	$3.7^{+3.0+1.3}_{-4.7-2.0}$	$3.7^{+2.9+1.0}_{-4.5-1.8}$	$-1.9^{+4.0}_{-3.5}$	$-2.5^{+5.3+3.3}_{-4.4-2.2}$	$-2.5^{+4.9+0.5}_{-4.1-1.0}$	-34^{+19}_{-19}

TABLE III. Ratios of the branching fractions in the SM and the nonuniversal Z' model.

	S1			S2			
R_i	SM	Case I	Case II	SM	Case I	Case II	Expt
R_1	$0.60^{+0.00}_{-0.04}$	$0.35^{+0.03+1.50}_{-0.02-0.15}$	$0.62^{+0.00+0.02}_{-0.05-0.03}$	$0.48^{+0.01}_{-0.03}$	$0.26^{+0.02+1.09}_{-0.04-0.12}$	$0.48^{+0.01+0.00}_{-0.03-0.00}$	$0.35^{+0.18}_{-0.15}$
R_2	$0.40^{+0.04}_{-0.00}$	$0.70^{+0.04+0.33}_{-0.09-0.66}$	$0.39^{+0.05+0.02}_{-0.00-0.01}$	$0.52^{+0.03}_{-0.01}$	$0.80^{+0.09+0.26}_{-0.05-0.75}$	$0.52^{+0.04+0.01}_{-0.01-0.00}$	
R_3	$1.00^{+0.03}_{-0.02}$	$0.87^{+0.28+0.59}_{-0.05-0.10}$	$1.00^{+0.04+0.01}_{-0.02-0.01}$	$0.99^{+0.02}_{-0.03}$	$1.05^{+0.08+0.27}_{-0.10-0.21}$	$0.99^{+0.02+0.01}_{-0.04-0.00}$	$1.25^{+0.36}_{-0.29}$

alleviate the disparity of $B^- \to \bar{K}_0^{*0}(1430)\pi^-$, but the large asymmetries in the $\bar{B}^0 \to K_0^{*-}(1430)\pi^+$ and $\bar{B}^0 \to K_0^{*0}(1430)\pi^0$ cannot be explained yet. In S2, for $\bar{B}^0 \to K_0^{*-}(1430)\pi^+$ and $\bar{B}^0 \to K_0^{*0}(1430)\pi^0$, the signs of center values are same as those of data. Furthermore, the CP asymmetries of $B^- \to \bar{K}_0^{*0}(1430)\pi^-$ in the SM and Z' model are almost null, which are close to the upper limits of experiment. Considering the large uncertainties, the results of S2 in both SM and Z' models can accommodate the data, except for the unexpectedly large asymmetry of $\bar{B}^0 \to K_0^{*0}(1430)\pi^0$, which should be measured critically in future. However, as pointed out in Ref. [37], final state interaction may have important effects on the decay rates and their direct CP violations, especially for the latter. However, this is beyond the scope of the present work.

Let us now analyze the impact of Z' on the isospin symmetry breaking. To explore the deviation from the isospin limit, it is convenient to define the following three parameters:

$$R_1 \equiv \frac{BR(\bar{B}^0 \to \bar{K}_0^{*0}(1430)\pi^0)}{BR(\bar{B}^0 \to K_0^{*-}(1430)\pi^+)}, \tag{31}$$

$$R_2 \equiv \frac{BR(B^- \to K_0^{*-}(1430)\pi^0)}{BR(B^- \to \bar{K}_0^{*0}(1430)\pi^-)}, \tag{32}$$

$$R_3 \equiv \frac{\tau(B^0)}{\tau(B^-)}\frac{BR(B^- \to \bar{K}_0^{*0}(1430)\pi^-)}{BR(\bar{B}^0 \to K_0^{*-}(1430)\pi^+)}. \tag{33}$$

Because they are the ratios of the branching fractions, they should be less sensitive to the nonperturbative inputs than other observables discussed before, therefore it is more persuasive to test them in both theoretical and experimental sides. In the isospin limits, i.e. ignoring the electroweak penguins, R_1, R_2, and R_3 are equal to 0.5, 0.5, and 1.0, respectively. So, the deviations reflect the magnitudes of the electroweak penguins directly. The results of SM and the nonuniversal Z' model are listed in Table. III. In the SM, it appears that the deviations from the isospin limit are not large in both scenarios, which shows that the QCD penguins are dominant. For Case I of the Z' model, the new physics just revise the Wilson coefficients of electroweak penguin operators, which could break the isospin symmetry. So, the ratios will be changed remarkably in both scenarios, as shown in the table. In Figure. 1, we also

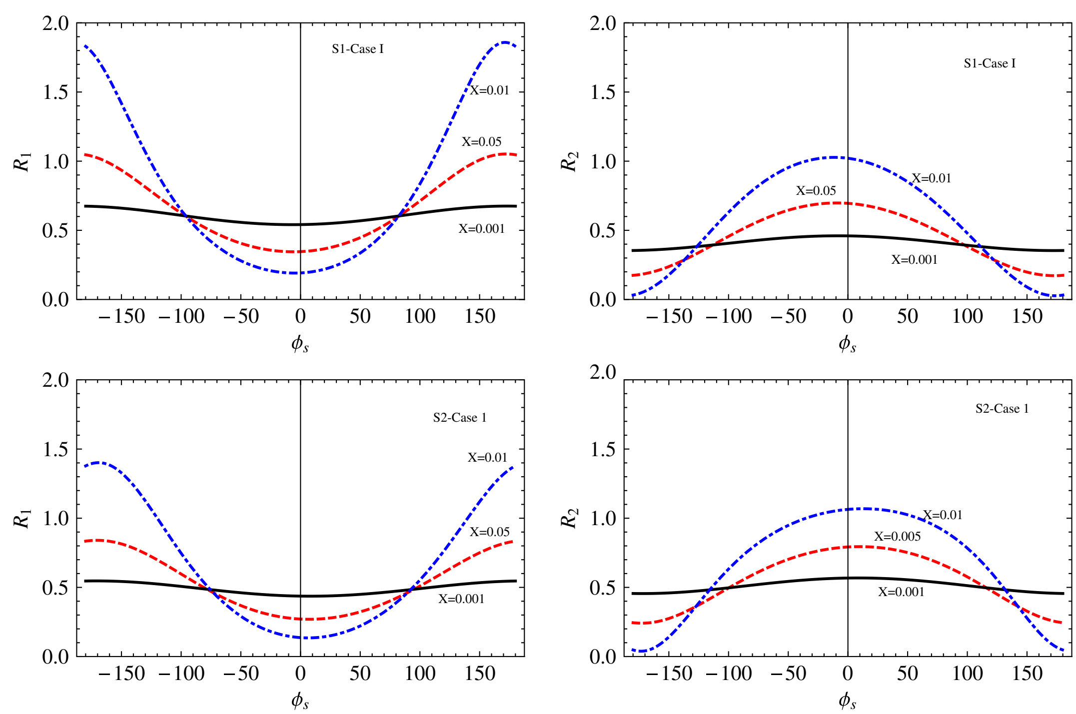

FIG. 1 (color online). R_1 and R_2 as functions of weak phase ϕ_s with different X in different scenarios and cases.

present the variations $R_{1,2}$ as functions of the new weak phase ϕ with different $X = 0.001$, 0.005, 0.01 in S1 (top panels) and S2 (bottom panels), so as to show the effect of two parameters X and ϕ. From the figures, we see that the $R_{1,2}$ change remarkably when $X = 0.01$ and 0.005. As $X = 0.001$, $R_{1,2}$ almost have same values as predictions of the SM. For Case II, the Z' boson changes the Wilson coefficients of QCD penguins, so the isospin symmetries are almost unchanged, as shown in Table. III. To sum up, the measurements of the R_i will help us determine whether QCD or electroweak interactions will be changed and then test the corresponding new physics models.

Finally, we will go back to the discussion of two scenarios. As aforementioned, $K_0^*(1430)$ is regarded as 2-quark state in both S1 and S2, but the only controversy is whether it belongs to ground state or the first excited state. Through calculation and comparison above, we favor the second scenario, which means that $K_0^*(1430)$ is the lowest lying $\bar{q}q$ state. Namely, the scalar mesons lower than 1 GeV are 4-quark states. This conclusion is also consistent with those of Refs. [3,12,15].

V. SUMMARY

Based on the QCD factorization approach, we have investigated in this work $B \to K_0^*(1430)\pi$ decays in the SM and a family nonuniversal Z' model. Because the inner structure of $K_0^*(1430)$ is not clear enough, we calculated the branching ratios under two different scenarios (S1 and S2). After calculation, we found that the branching ratios are sensitive to the weak annihilations. In the SM, with $\rho_A = 1$ and $\phi_A \in [-30°, 30°]$, the branching ratios of S1 (S2) are smaller (larger) than the experimental data. Considering the Z' boson in two different cases, for S1, the branching ratios are still far away from experiment. For S2, the branching ratios become smaller and can accommodate the data in Case II. In Case I, the results can also explain the data but with large uncertainties. Furthermore, the other interesting observables, such as CP asymmetries and isospin asymmetries, are also calculated. Compared with data, we favor that $K_0^*(1430)$ is the lowest lying $\bar{q}q$ state. Moreover, if there exists a Z' boson, Case II is preferable. All above results will be tested in the B factories, LHC-b, and the forthcoming super-B factory.

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[1] H. Y. Cheng and C. K. Chua, Nucl. Phys. B, Proc. Suppl. **207–208**, 391 (2010); A. J. Buras, Proc. Sci., BEAUTY (2011) 008.

[2] D. Asner *et al.* (Heavy Flavor Averaging Group), arXiv:1010.1589, and online update at http://www.slac.stanford.edu/xorg/hfag.

[3] H. Y. Cheng, C. K. Chua, and K. C. Yang, Phys. Rev. D **73**, 014017 (2006).

[4] M. Wirbel, B. Stech, and M. Bauer, Z. Phys. C **29**, 637 (1985); M. Bauer, B. Stech, and M. Wirbel, Z. Phys. C **34**, 103 (1987).

[5] A. Ali, G. Kramer, and C.-D. Lu, Phys. Rev. D **58**, 094009 (1998); **59**, 014005 (1998); H.-Y. Cheng and B. Tseng, Phys. Rev. D **58**, 094005 (1998).

[6] Y. Y. Keum, H. N. Li, and A. I. Sanda, Phys. Rev. D **63**, 054008 (2001); C. D. Lu, K. Ukai, and M. Z. Yang, Phys. Rev. D **63**, 074009 (2001).

[7] M. Beneke, G. Buchalla, M. Neubert, and C. T. Sachrajda, Phys. Rev. Lett. **83**, 1914 (1999); M. Beneke, G. Buchalla, M. Neubert, and C. T. Sachrajda, Nucl. Phys. **B591**, 313 (2000); M. Beneke and M. Neubert, Nucl. Phys. **B675**, 333 (2003).

[8] C. W. Bauer, D. Pirjol, and I. W. Stewart, Phys. Rev. Lett. **87**, 201806 (2001); Phys. Rev. D **65**, 054022 (2002).

[9] N. A. Tornqvist, Phys. Rev. Lett. **49**, 624 (1982); G. L. Jaffe, Phys. Rev. D **15**, 267 (1977); **15**, 281(E) (1977); A. L. Kataev, Yad. Fiz. **68**, 597 (2005); [Phys. At. Nucl. **68**, 567 (2005)]; J. Vijande, A. Valcarce, F. Fernandez, and B. Silvestre-Brac, Phys. Rev. D **72**, 034025 (2005).

[10] H. Y. Cheng and K. C. Yang, Phys. Rev. D **71**, 054020 (2005).

[11] H. Y. Cheng, C. K. Chua, and K. C. Yang, Phys. Rev. D **77**, 014034 (2008).

[12] C. H. Chen, Phys. Rev. D **67**, 014012 (2003).

[13] C. H. Chen, C. Q. Geng, Y. K. Hsiao, and Z. T. Wei, Phys. Rev. D **72**, 054011 (2005).

[14] W. Wang, Y. L. Shen, Y. Li, and C. D. Lu, Phys. Rev. D **74**, 114010 (2006).

[15] Y. L. Shen, W. Wang, J. Zhu, and C. D. Lu, Eur. Phys. J. C **50**, 877 (2007).

[16] C. s. Kim, Y. Li, and W. Wang, Phys. Rev. D **81**, 074014 (2010); X. Liu, Z. Q. Zhang, and Z.-J. Xiao, Chinese Phys. C **34**, 157 (2010); X. Liu and Z.-J. Xiao, Commun. Theor. Phys. **53**, 540 (2010); Z. Q. Zhang, Phys. Rev. D **82**, 034036 (2010); **82**, 114016 (2010); X. Liu, Z.-J. Xiao, and Z.-T. Zou, arXiv:1105.5761.

[17] C. H. Chen and C. Q. Geng, Phys. Rev. D **75**, 054010 (2007).

[18] S. Chaudhuri, S. W. Chung, G. Hockney, and J. Lykken, Nucl. Phys. **B456**, 89 (1995); M. Cvetic and P. Langacker, Phys. Rev. D **54**, 3570 (1996); G. Cleaver, M. Cvetic, J. R. Espinosa, L. L. Everett, P. Langacker, and J. Wang, Phys. Rev. D **59**, 055005 (1999); M. Cvetic, G. Shiu, and A. M. Uranga, Phys. Rev. Lett. **87**, 201801 (2001); M. Cvetic,

P. Langacker, and G. Shiu, Phys. Rev. D **66**, 066004 (2002).
[19] R. W. Robinett and J. L. Rosner, Phys. Rev. D **25**, 3036 (1982); **27**, 679(E) (1983).
[20] N. Arkani-Hamed, A. G. Cohen, and H. Georgi, Phys. Lett. B **513**, 232 (2001).
[21] P. Fayet, Phys. Rev. D **75**, 115017 (2007); C. H. Chen, C. Q. Geng, and C. W. Kao, Phys. Lett. B **663**, 400 (2008).
[22] P. Langacker and M.-x. Luo, Phys. Rev. D **45**, 278 (1992).
[23] S. Chaudhuri, S.-W. Chung, G. Hockney, and J. D. Lykken, Nucl. Phys. **B456**, 89 (1995).
[24] E. Nardi, Phys. Rev. D **48**, 1240 (1993); J. Bernabeu, E. Nardi, and D. Tommasini, Nucl. Phys. **B409**, 69 (1993); V. D. Barger, M. S. Berger, and R. J. Phillips, Phys. Rev. D **52**, 1663 (1995); M. B. Popovic and E. H. Simmons, Phys. Rev. D **62**, 035002 (2000); T. G. Rizzo, Phys. Rev. D **59**, 015020 (1998).
[25] P. Langacker and M. Plumacher, Phys. Rev. D **62**, 013006 (2000).
[26] P. Langacker, Rev. Mod. Phys. **81**, 1199 (2009).
[27] T. G. Rizzo, arXiv:hep-ph/0610104.
[28] A. Leike, Phys. Rep. **317**, 143 (1999).
[29] V. Barger, C.-W. Chiang, P. Langacker, and H.-S. Lee, Phys. Lett. B **580**, 186 (2004); **598**, 218 (2004); V. Barger, L. L. Everett, J. Jiang, P. Langacker, T. Liu, and C. E. M. Wagner, J. High Energy Phys. 12 (**2009**) 048; , Phys. Rev. D **80**, 055008 (2009).
[30] K. Cheung C.-W. Chiang, N. G. Deshpande, and J. Jiang, Phys. Lett. B **652**, 285 (2007); C.-W. Chiang N. G. Deshpande, and J. Jiang, J. High Energy Phys. 08 (2006) 075.
[31] Q. Chang, X. Q. Li, and Y. D. Yang, J. High Energy Phys. 05 (2009) 056; 02 (2010) 082; 04 (2010) 052; J. Hua, C. S. Kim, and Y. Li, Phys. Lett. B **690**, 508 (2010); Eur. Phys. J. C **69**, 139 (2010); Q. Chang and Y. H. Gao, Nucl. Phys. **B845**, 179 (2011); Y. Li, J. Hua, and K. C. Yang, Eur. Phys. J. C **71**, 1775 (2011).
[32] A. Arhrib, K. Cheung, C.-W. Chiang, and T.-C. Yuan, Phys. Rev. D **73**, 075015 (2006).
[33] C.-W. Chiang, Y.-F. Lin, and J. Tandean, J. High Energy Phys. 11 (2011) 083.
[34] C. D. Lu, Y. M. Wang, and H. Zou, Phys. Rev. D **75**, 056001 (2007).
[35] H. Y. Cheng, C. K. Chua, and C. W. Hwang, Phys. Rev. D **69**, 074025 (2004).
[36] G. Buchalla, A. J. Buras, and M. E. Lautenbacher, Rev. Mod. Phys. **68**, 1125 (1996); A. J. Buras, in *Probing the Standard Model of Particle Interactions*, Les Houches Lectures, edited by F. David and R. Gupta (Elsevier, New York, 1998).
[37] H. Y. Cheng, C. K. Chua, and A. Soni, Phys. Rev. D **71**, 014030 (2005).