Research Article

Branching Fractions and CP Asymmetries of $B \rightarrow K_0^*(1430)\rho(\omega)$ and $B \rightarrow K_0^*(1430)\phi$ Decays in the Family Nonuniversal Z' Model

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Within the QCD factorization framework, we investigate the branching fractions and the direct CP asymmetries of decays $B \rightarrow K_0^*(1430)\rho(\omega)$ and $B \rightarrow K_0^*(1430)\phi$ under two different scenarios in the standard model and the family nonuniversal Z' model. We find that the annihilation terms play crucial roles in these decays and lead to the major uncertainties. For decays $B^- \rightarrow K_0^{*-}(1430)\rho^0(\omega)$, the new Z' boson could change the branching fractions remarkably. However, for other decays, its contribution might be clouded by large uncertainties from annihilations. Unfortunately, neither the standard model nor the Z' model can reproduce all experimental data simultaneously under one certain scenario. We also noted that the direct CP asymmetries of $B^- \rightarrow K_0^{*-}(1430)\rho^0(\omega)$ could be used to identify the $K_0^*(1430)$ meson and search for the contribution of new Z' boson.

1. Introduction

The study of B meson rare decays is a crucial tool in testing the fundamental interactions among elementary particles, exploring the origin of CP violation, and searching for possible new physics (NP) beyond the standard model (SM). Theoretically and experimentally, such kind of research has been conducted in great detail, especially in the weak interactions of B meson. In particular, processes induced by flavor-changing neutral currents (FCNC) attract our attentions, because FCNC only occur at the loop level in SM, and they are always regarded as sensitive probes of NP. FCNC processes have been explored in the $B_q^0-\bar{B}_q^0$ mixing and the semileptonic weak decays, which permit a clean theoretical description. Charmless hadronic B meson decays, such as $B \rightarrow K\pi$, $K^{(*)}\phi$, and $K^{(*)}\eta^{(\prime)}$, induced by FCNC have been studied extensively. In the past few years, the effects of NP in these decays have been studied widely. These include supersymmetry models, two-Higgs doublet models, Z' models, fourth generation models, and extra dimension models (see review in [1] and references therein).

In the past twenty years, several novel methods have also been proposed to deal with the nonleptonic charmless B meson decays, such as the naive factorization [2, 3], the QCD

factorization (QCDF) [4–7], the perturbative QCD (PQCD) [8, 9], and the soft collinear effective theory [10–12]. In order to search for effect of NP in the hadronic B decays, the most theoretical studies are focused on $B \rightarrow PP$, PV , or VV . However, the studies of the decay modes involving a scalar meson are relatively few, because the underlying structure of the scalar mesons has not been well established theoretically. To describe the component of the scalar mesons, there are usually two possible scenarios (S1 and S2) according to the QCD sum rule method [13]: (i) in S1, we treat scalars above 1 GeV as the first excited states, while the scalars under 1 GeV are regarded as the low-lying states; (ii) in S2, the scalars above 1 GeV are viewed as the ground states, and light scalars are four-quark bound states or hybrid states. Under these two scenarios, many special decays have been examined within the QCDF [14, 15] and PQCD [16–25]. However, because of large uncertainties in SM, the NP effects in these decays are rarely studied.

Recently, the BaBar collaboration has reported their first branching fraction measurements for the decays $B \rightarrow K_0^*(1430)\rho$ that are induced by FCNC [26] and also updated their results of $B \rightarrow K_0^*(1430)\phi$ [27, 28] in [29]. Both experimental data and previous theoretical predictions are collected in Table 1 for comparison. For $B \rightarrow K_0^*(1430)\rho$,

TABLE 1: Branching ratios (in units of 10^{-6}) of $B \rightarrow K_0^*(1430)\rho$ and $B \rightarrow K_0^*(1430)\phi$ in QCDF [15], PQCD [19, 23], and experimental data [26, 29].

Decay modes	QCDF		PQCD		Data
	S1	S2	S1	S2	
$B^0 \rightarrow K_0^*(1430)\rho^0$	$10.0^{+2.4+0.5+12.1}_{-2.0-0.4-3.1}$	$36.0^{+13.8+0.9+23.2}_{-10.7-0.7-9.0}$	$0.47^{+0.12+0.20+0.03}_{-0.12-0.17-0.02}$	$4.8^{+1.1+1.0+0.3}_{-0.0-1.0-0.3}$	$27 \pm 4 \pm 2 \pm 3$
$B^0 \rightarrow K_0^*(1430)\rho^-$	$12.6^{+3.6+0.4+23.1}_{-3.0-0.3-5.8}$	$51.7^{+16.5+1.5+68.9}_{-13.4-1.4-24.0}$	$3.3^{+0.7+0.8+0.2}_{-0.6-0.8-0.2}$	$10.5^{+2.7+3.5+0.3}_{-0.0-2.6-0.3}$	$28 \pm 10 \pm 5 \pm 3$
$B^0 \rightarrow K_0^*(1430)\phi$	$0.9^{+0.3+0.4+19.3}_{-0.3-0.3-0.5}$	$16.9^{+6.2+1.7+51.8}_{-4.7-1.6-24.0}$	$3.7^{+0.8+0.1+3.7}_{-0.7-0.1-1.7}$	$23.6^{+5.6+0.8+10.9}_{-5.0-0.6-5.8}$	$3.9 \pm 0.5 \pm 0.6$
$B^\pm \rightarrow K_0^*(1430)\phi$	$1.0^{+0.3+0.4+20.2}_{-0.3-0.3-0.5}$	$17.3^{+6.2+1.7+52.4}_{-4.7-1.7-12.1}$	$4.3^{+0.9+0.1+4.3}_{-0.8-0.1-2.0}$	$25.6^{+6.2+0.9+12.1}_{-5.4-0.8-6.5}$	$7.0 \pm 1.3 \pm 0.9$

the data are inconsistent with the PQCD predictions in most cases. Moreover, these results are somewhat much lower (or larger) than the QCDF predictions for central values but are consistent with QCDF within rather large uncertainties. For $B \rightarrow K_0^*(1430)\phi$, all predicted central values deviate from the experimental data, though they can be also accommodated within very large theoretical errors. We note that in the following $K_0^*(1430)$ is denoted by K_0^* in some places for convenience.

The predictions of SM cannot agree with the data well, which permits us to search for possible new physics beyond SM in these decays. Our purpose of this work is to show that a new physics effect of similar size can be obtained from some models with an extra spin-1 Z' boson, which are known to naturally exist in some well-motivated extensions of SM [30]. Interesting phenomena arise when the Z' couplings to physical fermion eigenstates are nondiagonal, which could be realized in the E_6 models [31–35], string models [36], and some grand unified theories [37, 38]. For example, in the superstring model advocated by Chaudhuri et al. [36], it is possible to have family nonuniversal Z' couplings, because of different constructions of the different families. It also should be noted that, in such a model, called the family nonuniversal Z' model, the nonuniversal couplings could lead to FCNCs at the tree level as well as introduce new weak phases, which could explain the CP asymmetries in the current high energy experiments. In fact, the effects of Z' boson have been studied extensively in the low energy flavor physics phenomena, such as B meson mixing and decays [39–55], single top production [56], and lepton decays [57].

In this current work, we will adopt the QCDF approach [4–7] to evaluate the relevant hadronic matrix elements of B decays, since it is a systematic framework to calculate these matrix elements from QCD theory and holds in the heavy quark limit $m_b \rightarrow \infty$ and the heavy quark symmetry. In such calculations, one requires the additional knowledge about form factors of B meson to the scalar or the vector transitions. This problem, being a part of the nonperturbative sector of QCD, lacks a precise solution. To the best of our knowledge, a number of different approaches had been used to calculate the form factors of $B \rightarrow S$ decays, such as the QCD sum rule [58–60], light-cone QCD sum rule [61, 62], PQCD approach [63], and covariant light-front quark model (cLFQM) [64]. Among them, the form factors of the cLFQM are first calculated in the space-like region and their momentum dependence is fitted to a 3-parameter form. This parameterization is then analytically continued to the time-like region to determine

the physical form factors at $q^2 \geq 0$. Thus, we will use the results of cLFQM [64] in the following calculations.

We organize this paper as follows. In Section 2, we will reinvestigate $B \rightarrow K_0^*\rho$ and $K_0^*\phi$ in SM for comparison. In Section 3, we will review the family nonuniversal Z' model briefly and show the effect of Z' to decay modes we are considering. In Section 4, the numerical results and discussions are given. At last, this work will be summarized in Section 5.

2. Revisiting $B \rightarrow K_0^*(1430)\rho$ and $B \rightarrow K_0^*(1430)\phi$ Decays within the QCDF Framework

2.1. Input Parameters. Because most of parameters have been discussed and presented in [14, 15], we will just list the main inputs as follows.

2.1.1. Decay Constants. To proceed, we discuss the decay constants of the scalar meson. Unlike pseudoscalar meson, each scalar meson has two decay constants, the vector decay constant f_S and the scalar decay constant $\bar{f}_S$, namely, which are defined as

$$\langle S(p) | \bar{q}_2 \gamma_\mu q_1 | 0 \rangle = f_S p_\mu, \quad \langle S(p) | \bar{q}_2 q_1 | 0 \rangle = m_S \bar{f}_S, \quad (1)$$

and they are related by the equation of motion

$$f_S = \frac{m_2(\mu) - m_1(\mu)}{m_S} \bar{f}_S, \quad (2)$$

where m_2 and m_1 are the running current quark masses. Therefore, the vector decay constant is much smaller than the scalar one.

As for the vector meson, there are two kinds of decay constants, longitudinal decay constants and transverse decay constants, which are defined [65] as

$$\begin{aligned} \langle V(p) | \bar{q}_2 \gamma_\mu q_1 | 0 \rangle &= f_V m_V \epsilon_\mu^*, \\ \langle V(p, \epsilon^*) | \bar{q} \sigma_{\mu\nu} q' | 0 \rangle &= f_V^\perp (p_\mu \epsilon_\nu^* - p_\nu \epsilon_\mu^*). \end{aligned} \quad (3)$$

2.1.2. Distribution Amplitudes. In practice, the light-cone distribution amplitudes (LCDAs) of light mesons are required, which are nonperturbative and universal. In $B \rightarrow PP$,

PV, and VV decays [4–7], the twist-3 LCDAs are proven to be important and take about 30% contribution. Thus, we here use the LCDAs up to twist-3, and the discussions of higher twists can be found in [66]. The twist-2 $\phi_S(x)$ and twist-3 LCDAs of scalar mesons $\phi_S^s(x)$ and $\phi_S^\sigma(x)$ respect the normalization conditions

$$\int_0^1 dx \phi_S(x) = \frac{f_S}{2\sqrt{6}}, \quad (4)$$

$$\int_0^1 dx \phi_S^s(x) = \int_0^1 dx \phi_S^\sigma(x) = \frac{\bar{f}_S}{2\sqrt{6}},$$

and $\phi_S^T(x) = (1/6)(d/dx)\phi_S^\sigma(x)$. The twist-2 LCDA can be expanded in the Gegenbauer polynomials

$$\phi_S(x, \mu) = \frac{1}{\sqrt{6}} \bar{f}_S(\mu) 6x(1-x) \times \sum_{m=1}^{\infty} B_m(\mu) C_m^{3/2}(2x-1), \quad (5)$$

where B_m are Gegenbauer moments and $C_m^{3/2}$ are the Gegenbauer polynomials. The decay constants and the Gegenbauer moments of the twist-2 LCDA in two different scenarios have been studied explicitly in [14]. As for the explicit form of the Gegenbauer moments for the twist-3 LCDAs, there exist some uncertainties theoretically [67]; thus we adopt the asymptotic form for simplicity:

$$\phi_S^s = \frac{1}{\sqrt{6}} \bar{f}_S, \quad \phi_S^T = \frac{1}{\sqrt{6}} \bar{f}_S(1-2x). \quad (6)$$

For the vector mesons, the normalization for the twist-2 LCDA Φ_V and the twist-3 one Φ_v is given by

$$\int_0^1 dx \Phi_V(x) = f_V, \quad \int_0^1 dx \Phi_v(x) = 0, \quad (7)$$

where the definitions of $\Phi_v(x)$ can be found in [4–7]. The general expressions of these LCDAs read as

$$\begin{aligned} \Phi_V(x, \mu) &= 6x(1-x) f_V \left[1 + \sum_{n=1}^{\infty} \alpha_n^V(\mu) C_n^{3/2}(2x-1) \right], \\ \Phi_v(x, \mu) &= 3f_V^\perp \left[2x-1 + \sum_{n=1}^{\infty} \alpha_{n,\perp}^V(\mu) P_{n+1}(2x-1) \right], \end{aligned} \quad (8)$$

where $P_n(x)$ are the Legendre polynomials. The Gegenbauer moments α_n^V and $\alpha_{n,\perp}^V$ have been studied using the QCD sum rule method.

2.1.3. Form Factor. Another important nonperturbative parameters in our calculation are form factors of $B \rightarrow S, V$ transitions, which are defined by [2, 3]

$$\begin{aligned} \langle V(p') | V_\mu | B(p) \rangle &= -\frac{1}{m_B + m_V} \\ &\quad \times \varepsilon_{\mu\nu\alpha\beta} \varepsilon^{*\nu} P^\alpha q^\beta V^{BV}(q^2), \\ \langle V(p') | A_\mu | B(p) \rangle &= i \left\{ (m_B + m_V) \varepsilon_\mu^* A_1^{BV}(q^2) \right. \\ &\quad - \frac{\varepsilon^* \cdot P}{m_B + m_V} P_\mu A_2^{BV}(q^2) \\ &\quad \left. - 2m_V \frac{\varepsilon^* \cdot P}{q^2} q_\mu [A_3^{BV}(q^2) - A_0^{BV}(q^2)] \right\}, \\ \langle S(p') | A_\mu | B(p) \rangle &= -i \left[\left(P_\mu - \frac{m_B^2 - m_S^2}{q^2} q_\mu \right) F_1^{BS}(q^2) \right. \\ &\quad \left. + \frac{m_B^2 - m_S^2}{q^2} q_\mu F_0^{BS}(q^2) \right], \end{aligned} \quad (9)$$

with $P_\mu = (p + p')_\mu$, $q_\mu = (p - p')_\mu$. V and A denote the vector and axial-vector currents. As stated earlier, various form factors for $B \rightarrow S, V$ transitions have been evaluated in cLFQM [64], where form factors are first calculated in the space-like region and their momentum dependence is fitted to a 3-parameter form

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

The parameters a , b , and $F(0)$ relevant for our purposes are summarized in Table 2.

2.2. Amplitudes in QCD Factorization. We will use the QCD factorization approach to study the short-distance contribution of $B \rightarrow K_0^* \rho$ and $B \rightarrow K_0^* \phi$ decays. In SM, the effective weak Hamiltonian for $b \rightarrow s$ transitions is given by [68]

$$\begin{aligned} \mathcal{H}_{\text{eff}} &= \frac{G_F}{\sqrt{2}} \left[V_{ub} V_{us}^* (C_1 O_1^u + C_2 O_2^u) \right. \\ &\quad \left. + V_{cb} V_{cs}^* (C_1 O_1^c + C_2 O_2^c) - V_{tb} V_{ts}^* \right. \\ &\quad \left. \times \left(\sum_{i=3}^{10} C_i O_i + C_{7\gamma} O_{7\gamma} + C_{8g} O_{8g} \right) \right] + \text{h.c.}, \end{aligned} \quad (11)$$

where G_F is the Fermi coupling constant, $V_{qb} V_{qs}^*$ ($q = u, c, t$) are the products of Cabibbo-Kobayashi-Maskawa (CKM) matrix elements, O_i are the relevant four-quark operators whose explicit forms could be found, for example, in [68], and C_i are the corresponding Wilson coefficients.

TABLE 2: Form factors of $B \rightarrow \rho, K_0^*(1430)$ transitions obtained in the covariant light-front model [64].

F	$F(0)$	$F(q_{\max}^2)$	a	b
$V^{B\rho}$	0.27	0.79	1.84	1.28
$A_1^{B\rho}$	0.22	0.53	0.95	0.21
$F_1^{BK_0^*}$ [S1]	0.21	0.52	1.59	0.91
$F_1^{BK_0^*}$ [S2]	0.26	0.70	1.52	0.64
$A_0^{B\rho}$	0.28	0.76	1.73	1.20
$A_2^{B\rho}$	0.20	0.57	1.65	1.05
$F_0^{BK_0^*}$ [S1]	0.21	0.30	0.59	0.09
$F_0^{BK_0^*}$ [S2]	0.26	0.33	0.44	0.05

In the QCDF framework, the contribution of the matrix elements $\langle M_1 M_2 | O_i | B \rangle$ is dominated by the form factors and the nonfactorizable impact is controlled by hard gluon exchange. And the total elements can be written as

$$\begin{aligned} \langle M_1 M_2 | \mathcal{H}_{\text{eff}} | B \rangle \\ = \frac{G_F}{\sqrt{2}} \sum_{p=u,c} V_{pb} V_{ps}^* \langle M_1 M_2 | T_A^p + T_B^p | B \rangle, \end{aligned} \quad (12)$$

where T_A^p describes contributions from naive factorization, vertex corrections, penguin contractions, and spectator scattering expressed in terms of the flavor operators a_i^p , as shown in Figure 1. T_B^p contains annihilation topology amplitudes characterized by the annihilation operators b_j^p , as shown in Figure 2. In practice, the flavor operators a_i^p are basically the Wilson coefficients in conjunction with short-distance nonfactorizable corrections such as vertex corrections and hard spectator interactions. Combining the short-distance nonfactorizable corrections, the effective Wilson coefficients a_i^p have the expressions

$$\begin{aligned} a_i^p(M_1 M_2) = & \left(C_i + \frac{C_{i\pm 1}}{N_c} \right) N_i(M_2) \\ & + \frac{C_{i\pm 1}}{N_c} \frac{C_F \alpha_s}{4\pi} \\ & \times \left[V_i(M_2) + \frac{4\pi^2}{N_c} H_i(M_1 M_2) \right] \\ & + P_i^p(M_2), \end{aligned} \quad (13)$$

where $i = 1, \dots, 10$, the upper (lower) signs apply when i is odd (even), C_i are the Wilson coefficients, and $C_F = (N_c^2 - 1)/(2N_c)$ with N_c as the number of colors. The quantities $V_i(M_2)$ account for vertex corrections, $H_i(M_1 M_2)$ for hard spectator interactions with a hard gluon exchange between the emitted meson and the spectator quark of the B meson and $P_i(M_2)$ for penguin contractions. Weak annihilations are described by terms b_i and $b_{i,\text{EW}}$. For the explicit expressions of above functions, we refer the reader to [15] for details.

Equipped with these necessary preliminaries, the decay amplitudes could be expressed as

$$\begin{aligned} A(B^- \rightarrow K_0^{*-} \phi) \\ = i \frac{G_F}{\sqrt{2}} \sum_{p=u,c} \lambda_p^{(s)} \left\{ \left(a_3 + a_4^p + a_5 - r_\chi^\phi \left(a_6^p - \frac{1}{2} a_8^p \right) \right. \right. \\ \left. \left. - \frac{1}{2} (a_7 + a_9 + a_{10}^p) \right)_{K_0^* \phi} \right. \\ \left. \times 2 f_\phi F_1^{BK_0^*}(m_\phi^2) m_B p_c \right. \\ \left. - f_B f_\phi f_{K_0^*} (b_2 \delta_u^p + b_3 + b_{3,\text{EW}})_{K_0^* \phi} \right\}, \end{aligned}$$

$$\begin{aligned} A(\bar{B}^0 \rightarrow \bar{K}_0^{*0} \phi) \\ = i \frac{G_F}{\sqrt{2}} \\ \times \sum_{p=u,c} \lambda_p^{(s)} \left\{ \left(a_3 + a_4^p + a_5 - r_\chi^\phi \left(a_6^p - \frac{1}{2} a_8^p \right) \right. \right. \\ \left. \left. - \frac{1}{2} (a_7 + a_9 + a_{10}^p) \right)_{K_0^* \phi} \right. \\ \left. \times 2 f_\phi F_1^{BK_0^*}(m_\phi^2) m_B p_c \right. \\ \left. - f_B f_\phi f_{K_0^*} \left(b_3 - \frac{1}{2} b_{3,\text{EW}} \right)_{K_0^* \phi} \right\}, \end{aligned}$$

$$\begin{aligned} A(B^- \rightarrow \bar{K}_0^{*0} \rho^-) \\ = i \frac{G_F}{\sqrt{2}} \\ \times \sum_{p=u,c} \lambda_p^{(s)} \left\{ - \left(a_4^p + r_\chi^{K_0^*} \left(a_6^p - \frac{1}{2} a_8^p \right) - \frac{1}{2} a_{10}^p \right)_{\rho K_0^*} \right. \\ \left. \times 2 f_{K_0^*}^{B\rho} A_0^{B\rho}(m_{K_0^*}^2) m_B p_c \right. \\ \left. - f_B f_\rho f_{K_0^*} (b_2 \delta_u^p + b_3 + b_{3,\text{EW}})_{\rho K_0^*} \right\}, \end{aligned}$$

$$\begin{aligned} A(B^- \rightarrow K_0^{*-} \rho^0) \\ = i \frac{G_F}{2} \\ \times \sum_{p=u,c} \lambda_p^{(s)} \left\{ - \left(a_1 \delta_u^p + a_4^p + r_\chi^{K_0^*} (a_6^p + a_8^p) + a_{10}^p \right)_{\rho K_0^*} \right. \\ \left. \times 2 f_{K_0^*}^{B\rho} A_0^{B\rho}(m_{K_0^*}^2) m_B p_c \right. \\ \left. + \left[a_2 \delta_u^p + \frac{3}{2} (a_9 + a_7) \right]_{K_0^* \rho} \right. \\ \left. \times 2 f_\rho F_1^{BK_0^*}(m_\rho^2) m_B p_c \right. \\ \left. - f_B f_\rho f_{K_0^*} (b_2 \delta_u^p + b_3 + b_{3,\text{EW}})_{\rho K_0^*} \right\}, \end{aligned}$$

FIGURE 1: Order of α_s corrections to the hard scattering kernels. The upward quark lines represent the emission meson from the b quark decay vertex. These diagrams are commonly called vertex corrections, penguin corrections, and hard spectator scattering diagrams for (a)–(d), (e) and (f), and (g) and (h), respectively.

FIGURE 2: Order of α_s corrections to the weak annihilations.

$$A(\bar{B}^0 \rightarrow K_0^{*-} \rho^+)$$

$$= i \frac{G_F}{\sqrt{2}}$$

$$\times \sum_{p=u,c} \lambda_p^{(s)} \left\{ - \left(a_1 \delta_u^p + a_4^p + r_\chi^{K_0^*} a_6^p + a_{10}^p + r_\chi^{K_0^*} a_8^p \right)_{\rho K_0^*} \right. \\ \times 2 f_{K_0^*} A_0^{B\rho} (m_{K_0^*}^2) m_B p_c \\ \left. - f_B f_\rho f_{K_0^*} \left(b_3 - \frac{1}{2} b_{3,EW} \right)_{\rho K_0^*} \right\},$$

$$A(\bar{B}^0 \rightarrow \bar{K}_0^{*0} \rho^0)$$

$$= i \frac{G_F}{2}$$

$$\times \sum_{p=u,c} \lambda_p^{(s)} \left\{ - \left(-a_4^p - r_\chi^{K_0^*} \left(a_6^p - \frac{1}{2} a_8^p \right) + \frac{1}{2} a_{10}^p \right)_{\rho K_0^*} \right. \\ \times 2 f_{K_0^*} A_0^{B\rho} (m_{K_0^*}^2) m_B p_c \\ + \left[a_2 \delta_u^p + \frac{3}{2} (a_9 + a_7) \right]_{K_0^* \rho} \\ \times 2 f_\rho F_1^{BK_0^*} (m_\rho^2) m_B p_c \\ \left. - f_B f_\rho f_{K_0^*} \left(-b_3 + \frac{1}{2} b_{3,EW} \right)_{\rho K_0^*} \right\},$$

$$A(B^- \rightarrow K_0^{*-} \omega)$$

$$= i \frac{G_F}{2}$$

$$\times \sum_{p=u,c} \lambda_p^{(s)} \left\{ \left[a_2 \delta_u^p + 2(a_3 + a_5) + \frac{1}{2}(a_9 + a_7) \right]_{K_0^* \omega} \right. \\ \times 2 f_\omega F_1^{BK_0^*} (m_\omega^2) m_B p_c \\ - \left(a_1 \delta_u^p + a_4^p + r_\chi^{K_0^*} (a_6^p + a_8^p) + a_{10}^p \right)_{\omega K_0^*} \\ \times 2 f_{K_0^*} A_0^{B\omega} (m_{K_0^*}^2) m_B p_c \\ \left. - f_B f_\omega f_{K_0^*} (b_2 \delta_u^p + b_3 + b_{3,EW})_{\omega K_0^*} \right\},$$

$$A(\bar{B}^0 \rightarrow \bar{K}_0^{*0} \omega)$$

$$= i \frac{G_F}{2}$$

$$\times \sum_{p=u,c} \lambda_p^{(s)} \left\{ \left[a_2 \delta_u^p + 2(a_3 + a_5) + \frac{1}{2}(a_9 + a_7) \right]_{K_0^* \omega} \right. \\ \times 2 f_\omega F_1^{BK_0^*} (m_\omega^2) m_B p_c \\ - \left(a_4^p + r_\chi^{K_0^*} \left(a_6^p - \frac{1}{2} a_8^p \right) - \frac{1}{2} a_{10}^p \right)_{\omega K_0^*} \\ \times 2 f_{K_0^*} A_0^{B\rho} (m_{K_0^*}^2) m_B p_c \\ \left. - f_B f_\omega f_{K_0^*} \left(b_3 - \frac{1}{2} b_{3,EW} \right)_{\omega K_0^*} \right\},$$

where the ratios r_χ^V and r_χ^S are defined as

$$\begin{aligned} r_\chi^V(\mu) &= \frac{2m_V}{m_b(\mu)} \frac{f_V^\perp(\mu)}{f_V}, \\ r_\chi^S(\mu) &= \frac{2m_S^2}{m_b(\mu)(m_2(\mu) - m_1(\mu))}. \end{aligned} \quad (15)$$

The order of the arguments of the $a_i^P(M_1 M_2)$ and $b_i(M_1 M_2)$ coefficients is dictated by the subscript $M_1 M_2$, where M_1 shares the same spectator quark with the B meson and M_2 is the emitted meson. For the annihilation part, M_1 is referred to the one containing an antiquark from the weak vertex, and M_2 contains a quark from the weak vertex. And p_c is the c.m. momentum of the final mesons.

In QCDF, the endpoint singularities appear in calculating the twist-3 spectator and annihilation amplitudes. Since the treatment of endpoint divergences is model dependent, subleading power corrections generally can be studied only in a phenomenological way. As the most popular way, the endpoint divergent integrals are treated as signs of infrared sensitive contributions and parameterized by [4–7]

$$\int_0^1 \frac{dy}{y} \longrightarrow X_A = (1 + \rho_A e^{i\phi_A}) \ln \frac{m_B}{\Lambda_h}, \quad (16)$$

with the unknown real parameters ρ_A and ϕ_A . More discussion about them will be in Section 4.

3. The Family Nonuniversal Z' Model

In this section, we will review the main part of the family nonuniversal Z' model briefly. For simplicity, we only focus on the models in which the interactions between the Z' boson and fermions are flavor nonuniversal for left-handed couplings and flavor diagonal for right-handed cases. Of course, the analysis can be straightly extended to general cases in which the right-handed couplings are also nonuniversal across generations. The basic formulas of the Z' model with family nonuniversal and/or nondiagonal couplings have been presented in [30, 39], to which we refer readers for detail.

In the gauge basis, the neutral current Lagrangian induced by the Z' boson can be written as

$$\mathcal{L}^{Z'} = -g_2 J_\mu^{Z'} Z'^\mu, \quad (17)$$

where g_2 is the gauge coupling associated with the additional $U'(1)$ group at the M_W scale. Neglecting the renormalization group (RG) running effect between M_W and $M_{Z'}$ and the mixing between Z' and Z boson of SM, we present the chiral current as

$$J_\mu^{Z'} = \sum_{i,j} \bar{\psi}_i^I \gamma_\mu \left[(\varepsilon_{\psi_L})_{ij} P_L + (\varepsilon_{\psi_R})_{ij} P_R \right] \psi_j^I, \quad (18)$$

where the sum extends over the flavors of fermions, the chirality projection operators are $P_{L,R} \equiv (1 \mp \gamma_5)/2$, the superscript I stands for the weak interaction eigenstates, and $\varepsilon_{\psi_L}(\varepsilon_{\psi_R})$ denotes the left-handed (right-handed) chiral

coupling. ε_{ψ_L} and ε_{ψ_R} are required to be hermitian so as to arrive to a real Lagrangian. Accordingly, the mass eigenstates of the chiral fields can be defined by $\psi_{L,R} = V_{\psi_{L,R}} \psi_{L,R}^I$, and the usual CKM matrix is given by $V_{\text{CKM}} = V_{u_L} V_{d_L}^\dagger$. Then, the chiral Z' coupling matrices in the physical basis of up-type and down-type quarks are, respectively,

$$B_u^X \equiv V_{u_X} \varepsilon_{u_X} V_{u_X}^\dagger, \quad B_d^X \equiv V_{d_X} \varepsilon_{d_X} V_{d_X}^\dagger \quad (X = L, R). \quad (19)$$

If the ε matrices are not proportional to the identity, the B matrices will have nonzero off-diagonal elements, which induce FCNC interactions at the tree level directly. In this work, we assume that the right-handed couplings are diagonal for simplicity. Thereby, the effective Hamiltonian of the $\bar{b} \rightarrow \bar{s} q \bar{q}$ ($q = u, d$) transitions mediated by the Z' is

$$\begin{aligned} \mathcal{H}_{\text{eff}}^{Z'} &= \frac{2G_F}{\sqrt{2}} \left(\frac{g_2 M_Z}{g_1 M_{Z'}} \right)^2 B_{sb}^{L*} (\bar{b}s)_{V-A} \\ &\times \sum_q (B_{qq}^L (\bar{q}q)_{V-A} + B_{qq}^R (\bar{q}q)_{V+A}) + \text{h.c.}, \end{aligned} \quad (20)$$

where $g_1 = e/(\sin \theta_W \cos \theta_W)$ and $M_{Z'}$ is the mass of the new gauge boson. We note that the above operators of the forms $(\bar{b}s)_{V-A}(\bar{q}q)_{V-A}$ and $(\bar{b}s)_{V-A}(\bar{q}q)_{V+A}$ already exist in SM, so that we represent the Z' effect as a modification to the Wilson coefficients of the corresponding operators. Hence, we rewrite (20) as

$$\begin{aligned} \mathcal{H}_{\text{eff}}^{Z'} &= -\frac{G_F}{\sqrt{2}} V_{tb}^* V_{ts} \\ &\times \sum_q (\Delta C_3 O_3^{(q)} + \Delta C_5 O_5^{(q)} + \Delta C_7 O_7^{(q)} \\ &\quad + \Delta C_9 O_9^{(q)}) + \text{h.c.}, \end{aligned} \quad (21)$$

where the additional contributions to the SM Wilson coefficients at the M_W scale in terms of Z' parameters are given by

$$\begin{aligned} \Delta C_{3(5)} &= -\frac{2}{3V_{tb}^* V_{ts}} \left(\frac{g_2 M_Z}{g_1 M_{Z'}} \right)^2 B_{sb}^{L*} (B_{uu}^{L(R)} + 2B_{dd}^{L(R)}), \\ \Delta C_{9(7)} &= -\frac{4}{3V_{tb}^* V_{ts}} \left(\frac{g_2 M_Z}{g_1 M_{Z'}} \right)^2 B_{sb}^{L*} (B_{uu}^{L(R)} - B_{dd}^{L(R)}). \end{aligned} \quad (22)$$

Thus, we can have a Z' contribution to the QCD penguins $\Delta C_{3(5)}$ as well as the EW penguins $\Delta C_{9(7)}$, in the light of the results found by Buchalla et al. [68].

In order to simplify the calculations, many additional assumptions are often adopted, such as $B_{uu}^{L,R} = B_{dd}^{L,R}$ and $B_{uu}^{L,R} = -2B_{dd}^{L,R}$. In this paper, we adopt the latter one, because we hope the new physics is primarily manifest in the EW penguins and violates the isospin symmetries. In the literatures, this assumption has been used widely [42, 43, 46–50, 56]. As a

result, the Z' contributions to the Wilson coefficients at the weak scale are

$$\begin{aligned}\Delta C_{3(5)} &= 0, \\ \Delta C_{9(7)} &= 4 \frac{|V_{tb}^* V_{ts}|}{V_{tb}^* V_{ts}} \xi^{LL(R)} e^{-i\phi_L},\end{aligned}\quad (23)$$

where

$$\begin{aligned}\xi^{LX} &\equiv \left(\frac{g_2 M_Z}{g_1 M_{Z'}} \right)^2 \left| \frac{B_{sb}^{L*} B_{dd}^X}{V_{tb}^* V_{ts}} \right| \quad (X = L, R), \\ \phi_L &\equiv \text{Arg} [B_{sb}^L].\end{aligned}\quad (24)$$

Because of the hermiticity of the effective Hamiltonian, the diagonal elements of the effective coupling matrix must be real. However, the off-diagonal elements, such as B_{sb}^L , generally may contain new weak phases. Moreover, the relation $B_{ss}^{L(R)} \simeq B_{dd}^{L(R)}$ follows from the assumptions of universality for the first two families, as required by K and μ decay constraints [39].

It should be emphasized that the other SM Wilson coefficients may also receive contributions from the Z' boson through RG evolution. With our assumption that there is 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 SM [68]. The numerical results of Wilson coefficients in the naive dimensional regularization (NDR) scheme at the scale $\mu = 2.1 \text{ GeV}$ ($\mu_h = 1 \text{ GeV}$) are listed in Table 3 for convenience.

In summary, we list here our simplifications to a general Z' model: we assume (i) no right-handed flavor-changing couplings ($B_{ij}^R = 0$ for $i \neq j$), (ii) no significant RG running effect between $M_{Z'}$ and M_W scales, (iii) negligible Z' effect on the QCD penguin ($\Delta C_{3,5} = 0$) so that the new physics is manifestly isospin violating. With these simplifications, we have only three parameters left in the model, $|\xi_{qq}^L|$, $|\xi_{qq}^R|$, and ϕ_L . So, this approach provides a minimal way to introduce the Z' effect in the concerned decay modes. Of course, more general Z' models are possible.

Now, the only task left is to constrain the parameters within the existing experimental data. Generally, $g_2/g_1 \sim 1$ is expected, if both the $U(1)$ gauge groups have the same origin from some grand unified theories. We also hope $M_{Z'}/M_W \sim 0.1$ so that TeV scale neutral Z' boson could be detected at LHC. Theoretically, one can fit the left three parameters $|B_{sb}^L|$, $|B_{dd}^X|$, and new weak phase ϕ_L with the accurate data from B factories and other experiments such as Tavatron and LHC. For example, B_{sb}^L and ϕ_L could be extracted from $B_s - \bar{B}_s$ mixing as well as $B \rightarrow K^{(*)} \ell^+ \ell^-$ decays. To accommodate the mass difference between B_s and $\bar{B}_s$, $|B_{sb}^L| \sim |V_{tb} V_{ts}^*|$ is required [42–44, 48–50]. In [48–50], the authors got that the ϕ_L is about -80° by fitting data of $B_s - \bar{B}_s$ mixing and $B \rightarrow K^{(*)} l^+ l^-$ decays. Subsequently, with B_{sb}^L and ϕ_L arrived and experimental data of $B \rightarrow \pi\pi, K\pi, K\rho$, and $K^{(*)}\phi$, B_{qq}^L and B_{qq}^R could be extracted analogously. Specifically, the CP asymmetries of $B \rightarrow K\pi$ and polarizations of

$B \rightarrow K^* \phi$ constrain $|B_{sb}^L B_{ss}^{LR}| \sim |V_{tb} V_{ts}^*|$, which indicates $|B_{qq}^{L,R}| \sim 1$. However, we have one remark here: in dealing with the nonleptonic B decays, because different groups used different factorization approaches, the fitted results are different, but all results have a same order. Note that the detailed constraint of these parameters is beyond the scope of current work and can be found in many references [46–50]. Summing up above analysis, we thereby assume that $\xi^{LL(R)} \in (10^{-3}, 10^{-2})$ and $\phi_L \in (-60^\circ, -90^\circ)$ so as to probe the new physics effect for a maximum range.

4. Numerical Results and Discussion

We will discuss the decay constants of scalars, firstly. Since it is not clear whether the scalar meson $K_0^*(1430)$ belongs to the first orbital excited states (S1) or the low-lying resonances (S2), we will calculate the processes under both scenarios. In the calculation, the decay constants and Gegenbauer moments obtained within the QCD sum rules method under different scenarios are presented as follows [14]:

$$\begin{aligned}\text{S1: } \bar{f}_{K_0^*}(1.0 \text{ GeV}) &= -(300 \pm 30) \text{ MeV}; \\ \bar{f}_{K_0^*}(2.1 \text{ GeV}) &= -(370 \pm 35) \text{ MeV}; \\ B_1(1.0 \text{ GeV}) &= 0.58 \pm 0.07; \\ B_1(2.1 \text{ GeV}) &= 0.39 \pm 0.05; \\ B_3(1.0 \text{ GeV}) &= -1.20 \pm 0.08; \\ B_3(2.1 \text{ GeV}) &= -0.70 \pm 0.05; \\ \text{S2: } \bar{f}_{K_0^*}(1.0 \text{ GeV}) &= (445 \pm 50) \text{ MeV}; \\ \bar{f}_{K_0^*}(2.1 \text{ GeV}) &= (550 \pm 60) \text{ MeV}; \\ B_1(1.0 \text{ GeV}) &= -0.57 \pm 0.13; \\ B_1(2.1 \text{ GeV}) &= -0.39 \pm 0.09; \\ B_3(1.0 \text{ GeV}) &= -0.42 \pm 0.22; \\ B_3(2.1 \text{ GeV}) &= -0.25 \pm 0.13.\end{aligned}\quad (25)$$

The longitudinal and transverse decay constants of vector mesons are listed as

$$\begin{aligned}f_\rho &= 216 \text{ MeV}, & f_\omega &= 187 \text{ MeV}, & f_\phi &= 215 \text{ MeV}, \\ f_\rho^\perp &= 165 \text{ MeV}, & f_\omega^\perp &= 151 \text{ MeV}, & f_\phi^\perp &= 186 \text{ MeV},\end{aligned}\quad (26)$$

where the values are taken from [69]. In the LCADs of vectors, the Gegenbauer moments α_n^V and $\alpha_{n,\perp}^V$ have been studied, and we will employ the most recent updated values from QCD sum rules [70]:

$$\begin{aligned}\alpha_2^{\rho,\omega} &= 0.15, & \alpha_{2,\perp}^{\rho,\omega} &= 0.14, \\ \alpha_2^\phi &= 0.18, & \alpha_{2,\perp}^\phi &= 0.14,\end{aligned}\quad (27)$$

and $\alpha_1^V = 0, \alpha_{1,\perp}^V = 0$.

TABLE 3: The Wilson coefficients C_i within SM and with the contribution from Z' boson included in NDR scheme at the scale $\mu = 2.1$ GeV and $\mu_h = 1.0$ GeV.

Wilson coefficients	$\mu = 2.1$ GeV		$\mu_h = 1.0$ GeV	
	C_i^{SM}	$\Delta C_i^{Z'}$	C_i^{SM}	$\Delta C_i^{Z'}$
C_1	1.135	0	1.224	0
C_2	-0.283	0	-0.429	0
C_3	0.021	$0.09\xi^{LL} - 0.02\xi^{LR}$	0.034	$0.15\xi^{LL} - 0.04\xi^{LR}$
C_4	-0.049	$-0.20\xi^{LL} + 0.01\xi^{LR}$	-0.072	$-0.31\xi^{LL} + 0.03\xi^{LR}$
C_5	0.010	$0.03\xi^{LL} + 0.02\xi^{LR}$	0.010	$0.02\xi^{LL} + 0.02\xi^{LR}$
C_6	-0.06	$-0.26\xi^{LL} + 0.03\xi^{LR}$	-0.104	$-0.44\xi^{LL} + 0.07\xi^{LR}$
C_7/α_{em}	-0.018	$5.3\xi^{LL} - 461\xi^{LR}$	-0.023	$6.3\xi^{LL} - 457\xi^{LR}$
C_8/α_{em}	0.081	$2.43\xi^{LL} - 286\xi^{LR}$	0.134	$4.8\xi^{LL} - 497\xi^{LR}$
C_9/α_{em}	-1.266	$-594\xi^{LL} + 6.1\xi^{LR}$	-1.366	$-643\xi^{LL} + 7.8\xi^{LR}$
$C_{10}/\alpha_{\text{em}}$	0.321	$178\xi^{LL} - 1.0\xi^{LR}$	0.483	$257\xi^{LL} - 1.9\xi^{LR}$
$C_{7\gamma}$	-0.345	—	-0.395	—
C_{8g}	-0.161	—	-0.181	—

In [14, 15, 25], it was found that in decay modes with scalars the main theoretical uncertainties are due to the annihilations, especially for the penguin dominated ones. In $B \rightarrow PP, PV$ decays, the annihilation amplitudes are helicity suppressed because the helicity of one of the final states cannot match with that of its quarks. However, this helicity suppression can be alleviated in the decay modes with scalar because of nonvanishing orbital angular momentum. Thus, annihilation contribution to $B \rightarrow SP(V)$ is much larger than the $B \rightarrow PP(V)$ case. However, as stated before, the end-point singularity appears in calculating the annihilation contribution, and then two free parameters, ρ_A and ϕ_A , are introduced phenomenally. In [15], it is found that the behavior of SV is similar to the longitudinal part of VV . Fortunately, with experimental data, it presents the moderate value of nonuniversal annihilation phase $\phi_A = -40^\circ$ for $B \rightarrow VV$ decay modes [4–7]. Therefore, for $B \rightarrow SV$, we conservatively take $\phi_A = (-40 \pm 20)^\circ$ with $\rho_A = 0.6 \pm 0.2$, which also assures that the hadronic uncertainties are considerably reduced. Furthermore, the endpoint divergence X_H in the hard spectator contributions can also be parameterized in the same manner.

Within above parameters and formulas, we calculate the branching fractions of these decays in SM and the family nonuniversal Z' model under two different scenarios. Together with partial experimental results, the results are exhibited in Table 4, respectively. For the center values, we adopt $\xi^{LL,LR} = 0.005$ and $\phi_L = -80^\circ$. For all theoretical predictions, the first errors involved the uncertainties of the Gegenbauer moments $B_{1,3}$, the scalar meson decay constants, and form factors. The second errors arise from the power corrections of annihilation and hard spectator interactions characterized by the parameters $X_{A,H}$. To obtain the third errors of the Z' model results, we scan randomly the points in their own possible parameter spaces.

Compared with predictions in [15] (considering the typos), our results agree with theirs well within uncertainties. However, there still exist few differences, and some reasons

are list as follows: (1) in [15], for the parameterizations of singularities, the center values correspond to $\rho_{A,H} = 0$ and $\phi_{A,H} = 0$, while we set $\rho_{A,H} = 0.6$ and $\phi_{A,H} = -40^\circ$; (2) the difference of Wilson coefficients, caused by the top quark mass and other part parameters, will change the results slightly; (3) in this work, the different form factors of $B \rightarrow K_0^*(1430)$ are used under different scenarios, but they adopted same values in [15].

In Table 4, for $K_0^*(1430)\phi$ channels, though the central values of the prediction under S1 are smaller than the experimental data, they are accommodated with the large uncertainties. However, in S2, the theoretical results are much larger than the data and cannot agree with data even with uncertainties. These theoretical results also agree with the predictions of PQCD [19]. For $\bar{B}^0 \rightarrow \bar{K}_0^{*0}\rho^0$ channels, contrary to $K_0^*(1430)\phi$, the result of S2 agrees with data well and the prediction of S1 is much smaller than the data. Since for $\bar{B}^0 \rightarrow K_0^{*-}\rho^+$ there is large uncertainty in the experimental data, the theoretical results under both S1 and S2 can accommodate the data with large uncertainties theoretically. That is to say, it is impossible to explain all data under one settled scenario simultaneously.

When adding the contribution of the Z' gauge boson, as shown in the table, the Z' gauge boson changes the branching fractions under both two different scenarios. For $B \rightarrow K_0^*\phi$ channels dominated by the annihilation, the Z' will enhance the branching fractions in S1, while in S2 the branching ratios are decreased. The reason is that the annihilation is proportional to the decay constant $f_{K_0^*}$, which has different sign in different scenarios. For $B^- \rightarrow \bar{K}_0^{*0}\rho^-$ and $\bar{B}^0 \rightarrow K_0^{*-}\rho^+$, because the scalar particle is the emitted particle, the whole amplitudes are proportional to the decay constant $f_{K_0^*}$; thus the new physics contribution has same trends in different scenarios. For channels with ρ^0 or ω , the spectator quarks enter not only the scalars but also the vectors, the amplitudes become much complicated, and we cannot describe the relations between new physics and branching fractions easily.

TABLE 4: Branching fractions (in units of 10^{-6}) under the different scenarios.

Decay mode	S1		S2		Expt.
	SM	SM + Z'	SM	SM + Z'	
$B^- \rightarrow K_0^{*-} \phi$	$2.4^{+0.2+5.1}_{-0.2-1.8}$	$3.8^{+0.2+5.6+1.8}_{-0.2-2.5-1.2}$	$22.6^{+2.0+19.7}_{-4.5-8.6}$	$16.5^{+1.2+18.5+5.0}_{-4.1-10.3-7.9}$	$7.0 \pm 1.3 \pm 0.9$
$\bar{B}^0 \rightarrow K_0^{*0} \phi$	$2.2^{+0.2+4.9}_{-0.2-1.7}$	$4.7^{+0.4+5.2+4.4}_{-0.4-2.6-2.2}$	$22.4^{+2.0+19.4}_{-4.4-8.4}$	$21.2^{+2.2+18.9+3.0}_{-4.1-8.4-4.6}$	$4.3 \pm 0.6 \pm 0.4$
$B^- \rightarrow \bar{K}_0^{*0} \rho^-$	$11.7^{+2.5+8.4}_{-2.2-4.0}$	$11.5^{+2.4+7.7+3.6}_{-2.2-4.7-3.8}$	$45.5^{+9.7+20.6}_{-9.6-10.4}$	$41.4^{+8.9+19.7+4.4}_{-8.7-11.9-5.7}$	
$B^- \rightarrow K_0^{*-} \rho^0$	$7.2^{+1.2+4.5}_{-1.3-2.3}$	$18.2^{+2.3+5.9+27.4}_{-3.2-3.4-10.0}$	$17.6^{+5.3+9.0}_{-4.4-4.4}$	$15.9^{+5.2+8.4+6.9}_{-4.2-6.1-5.6}$	
$\bar{B}^0 \rightarrow \bar{K}_0^{*0} \rho^0$	$4.6^{+1.1+1.7}_{-1.0-0.7}$	$3.9^{+0.8+1.4+6.4}_{-0.6-0.8-1.8}$	$24.5^{+4.8+7.0}_{-5.4-3.8}$	$33.3^{+5.0+7.7+32.1}_{-5.7-4.8-8.5}$	27 ± 5.5
$\bar{B}^0 \rightarrow K_0^{*-} \rho^+$	$10.7^{+2.3+8.5}_{-2.0-3.7}$	$14.4^{+3.0+9.2+8.7}_{-2.7-5.1-4.1}$	$44.7^{+10.6+21.3}_{-9.5-10.5}$	$54.1^{+12.9+22.4+25.7}_{-11.4-12.6-8.6}$	28 ± 12
$B^- \rightarrow K_0^{*-} \omega$	$3.6^{+0.7+3.1}_{-0.7-1.4}$	$7.8^{+1.4+3.7+10.1}_{-1.3-1.7-3.8}$	$12.6^{+3.0+7.2}_{-2.7-3.5}$	$13.7^{+4.0+7.5+4.8}_{-2.8-3.9-1.4}$	
$\bar{B}^0 \rightarrow \bar{K}_0^{*0} \omega$	$3.9^{+0.8+1.3}_{-0.7-0.6}$	$4.0^{+0.8+1.3+1.2}_{-0.7-0.7-0.6}$	$10.6^{+2.7+4.4}_{-2.2-2.1}$	$10.7^{+2.7+4.4+5.4}_{-2.2-2.7-2.6}$	

Compared to the experimental data, the Z' boson affects the branching fractions remarkably and alleviates the disparities. However, we cannot achieve a definite conclusion yet whether K_0^* belongs to the ground states or the first orbital excited states. Moreover, for most modes except $B^- \rightarrow K_0^{*-} \rho^0(\omega)$, the new physics contribution might be clouded by the uncertainties taken by the annihilations. Thus, it is also very difficult to search for Z' effect in these decays. Specifically, for decays $B^- \rightarrow K_0^{*-} \rho^0(\omega)$, Z' boson could enhance the branching fractions more than two times; we hope these two channels could be measured in the LHCb experiment or Super-b factories in future so as to probe the Z' gauge boson.

To test the isospin symmetry and probe new physics, we define two ratios:

$$R_1 = \frac{Br(\bar{B}^0 \rightarrow \bar{K}_0^{*0} \rho^0)}{Br(\bar{B}^0 \rightarrow K_0^{*-} \rho^+)} = 0.96^{+1.07}_{-0.43} [\text{Exp.}], \quad (28)$$

$$R_2 = \frac{\tau(\bar{B}^0)}{\tau(B^-)} \cdot \frac{Br(B^- \rightarrow K_0^{*-} \phi)}{Br(\bar{B}^0 \rightarrow \bar{K}_0^{*0} \phi)} = 1.68^{+0.90}_{-0.60} [\text{Exp.}],$$

where the experimental results are also given, and all uncertainties are added in quadrature. In the isospin limit, $R_1 = 1/2$ and $R_2 = 1$ are expected to hold. Here, we list the theoretical results under different scenarios in different models:

$$\begin{aligned}
R_1 [\text{SM}] &= \begin{cases} 0.43^{+0.13}_{-0.10}, & \text{S1} \\ 0.55^{+0.08}_{-0.08}, & \text{S2}; \end{cases} \\
R_1 [\text{SM} + Z'] &= \begin{cases} 0.27^{+0.15+0.31}_{-0.08-0.17}, & \text{S1} \\ 0.62^{+0.09+0.20}_{-0.09-0.07}, & \text{S2}; \end{cases} \\
R_2 [\text{SM}] &= \begin{cases} 1.00^{+0.04}_{-0.04}, & \text{S1} \\ 0.94^{+0.00}_{-0.01}, & \text{S2}; \end{cases} \\
R_2 [\text{SM} + Z'] &= \begin{cases} 0.76^{+0.16+0.24}_{-0.18-0.18}, & \text{S1} \\ 0.73^{+0.16+0.19}_{-0.27-0.24}, & \text{S2}. \end{cases}
\end{aligned} \quad (29)$$

In the above results, the theoretical uncertainties are reduced since they are ratios of branch fractions. Because R -values

are ratios of branching fractions, the uncertainties taken by nonperturbative part are almost zero, and we here ignored them. We see that the symmetries are almost held in SM, while the data shows that the isospin symmetries are violated, which means that the large electroweak penguins or weak annihilations may break the isospin symmetries remarkably. When adding Z' contribution, except for R_1 under S2, the isospin symmetries are broken in opposite directions. However, the family nonuniversal Z' model cannot be ruled out due to large uncertainties in both experimental and theoretical sides.

Finally, we will discuss the direct CP asymmetries of these decays. Because $|V_{ub}V_{us}|(\lambda^4) \ll |V_{tb}V_{ts}|(\lambda^2)$ for the charged mode $B^- \rightarrow K_0^{*-} \phi$ and there is no tree contribution in the neutral mode, thereby the direct CP asymmetries of $B \rightarrow K_0^* \phi$ are almost zero in both SM and the Z' model. For $B \rightarrow K_0^* \rho$, although the CKM elements are suppressed, the tree operators with large Wilson coefficients appear in the emission diagrams, so the amplitudes of tree and penguin may have comparable magnitudes. Thus, large CP asymmetries in these decays are expected, just like decays $B \rightarrow K\pi$ and $B \rightarrow K\rho$. In Table 5, we list the direct CP asymmetries of $B \rightarrow K_0^* \rho$ under different scenarios. In this table, the uncertainties induced by nonperturbative part are not included, as the direct CP asymmetries are also ratios. From the table, we firstly note that $\bar{B}^0 \rightarrow \bar{K}_0^{*0} \rho(\omega)$ have large asymmetries, and different scenarios have different signs but with large uncertainties. If we can calculate the annihilation accurately within some effective approach in future, this parameter could be used to distinguish the scenarios. Secondly, for $B^- \rightarrow K_0^{*-} \rho^0(\omega)$, the Z' could change the signs of the center values, and these two decays can be used in probing new physics effect.

5. Summary

Motivated by recent measurements of decays $B \rightarrow K_0^* \rho$ and $K_0^* \phi$, we studied the branching fractions of these decays both in SM and in the family nonuniversal Z' model within the QCDF framework. Because it is not clear whether K_0^* is the lying state or the first orbital excited state, we calculated them under two different scenarios. We noted that, for these decay modes with scalar meson, the annihilations

TABLE 5: The direct CP asymmetry (%) under the different scenarios.

Decay mode	S1		S2	
	SM	SM + Z'	SM	SM + Z'
$B^- \rightarrow \bar{K}_0^{*0} \rho^-$	6_{-2}^{+4}	6_{-3-1}^{+6+1}	2_{-1}^{+2}	2_{-1-0}^{+2+0}
$B^- \rightarrow K_0^{*-} \rho^0$	4_{-3}^{+4}	-3_{-1-2}^{+2+4}	-1_{-4}^{+3}	6_{-3-6}^{+3+10}
$\bar{B}^0 \rightarrow \bar{K}_0^{*0} \rho^0$	9_{-38}^{+26}	24_{-42-12}^{+52+21}	-11_{-13}^{+10}	-9_{-11-2}^{+9+4}
$\bar{B}^0 \rightarrow K_0^{*-} \rho^+$	1_{-2}^{+1}	-2_{-1-0}^{+1+1}	1_{-0}^{+0}	1_{-0-0}^{+0+0}
$B^- \rightarrow K_0^{*-} \omega$	3_{-7}^{+6}	-4_{-3-3}^{+2+5}	-1_{-5}^{+4}	-4_{-4-6}^{+3+5}
$\bar{B}^0 \rightarrow \bar{K}_0^{*0} \omega$	16_{-39}^{+26}	17_{-40-5}^{+28+7}	-19_{-15}^{+14}	-4_{-19-13}^{+15+19}

play more important roles than those in $B \rightarrow PP$ and PV decays; thus annihilations lead to large uncertainties in these decays. From this point, an effective way that could calculate the annihilations reliably is needed. Compared with the experimental results, we found that different decay modes favor different scenarios. Moreover, in order to account for the large isospin asymmetries in the data, large weak annihilations are also required. Adding the contribution of the family nonuniversal Z' boson, the branching fractions and their ratios are both changed remarkably. However, we cannot identify the character of the scalar meson K_0^* , either. Furthermore, for most channels, the Z' contribution will be buried by large uncertainties, except for decays $B^- \rightarrow K_0^{*-} \rho^0(\omega)$. For the isospin asymmetries, the Z' boson makes the disparities more worse, but we cannot rule this model out because of large errors in both experimental and theoretical sides.

In this work, we also calculated the direct CP asymmetries and found that the CP asymmetries of $B \rightarrow K_0^* \phi$ are almost zero. In different scenarios, the CP asymmetries of $B^- \rightarrow K_0^{*-} \rho^0(\omega)$ have different signs; as a result, they can be used to classify the scalar K_0^* . If its character could be identified, we accordingly could use these results to probe the new gauge boson Z' , because it also could changes the signs of the direct CP asymmetries. All the above predictions could be tested in the running LHCb or the Super-b factories in the future.

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