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CP violation in $B \to \phi K$ decay with anomalous right-handed top quark couplings

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Abstract. We explore CP violation in $B \to \phi K$ decay processes in the presence of the anomalous right-handed $\bar{t}sW$ and $\bar{t}bW$ couplings. The complex anomalous top coupling can be a source of new CP violation and may lead to a deviation of the observed weak phase in $B \to \phi K$ decays, which accounts for the present disagreement of the observed $\sin 2\beta$ between $B \to J/\psi K$ and $B \to \phi K$ decays. Direct CP violation is also predicted.

1 Introduction

Recently the BaBar [1] and Belle [2] Collaborations reported the first measurement of the time-dependent CP asymmetry in $B \to \phi K$ decay to measure the weak phase $\sin 2\beta$:

$$\begin{aligned} \sin 2\beta &= -0.73 \pm 0.64 \pm 0.18 \quad \text{(Belle)}, \\ \sin 2\beta &= -0.19^{+0.52}_{-0.50} \pm 0.09 \quad \text{(BaBar)}, \end{aligned} \tag{1}$$

where $\beta \equiv \arg(-V_{cd}V_{cb}^*/V_{td}V_{tb}^*)$. In the standard model (SM), the origin of CP violation is only the complex phase of the Cabibbo–Kobayashi–Maskawa (CKM) quark mixing matrix. It implies that $\sin 2\beta$ in $B \to \phi K$ decays should agree with that of $B \to J/\psi K$ decays up to small pollution of $\mathcal{O}(\lambda^2) \sim 5\%$ [3,4]. Therefore a sizable disagreement of $\sin 2\beta$ between the $B \to \phi K$ and $B \to J/\psi K$ decays is a clear indication of new physics beyond the SM. The world average of $\sin 2\beta$ measured in $B \to J/\psi K$ decays is given by [5]

$$\sin 2\beta = 0.734 \pm 0.054, \tag{2}$$

which is consistent with the SM prediction and indicates the non-zero CP violation in the B system. Remarkably, however, the measured $\sin 2\beta$ in the $B \to \phi K_S$ channel is far from that of $B \to J/\psi K$ decay and even the central value is negative as shown in (1). At present, we are confronted with a 2.7σ discrepancy between the average values of $\sin 2\beta$ in $B \to \phi K_S$ and in $B \to J/\psi K_S$ decays. Although it is premature to regard this disagreement as evidence of new physics due to the large statistical error, the difference is so large that it is tempting to interpret it as a clue of new physics. Studies in various models are being performed to account for the discrepancy [5,6].

The left–right (LR) model based on the $SU(2)_L \times SU(2)_R \times U(1)$ gauge group is one of the natural extensions of the SM [7]. In the LR model, right-handed quark mixing is also an observable as well as left-handed quark mixing. Without a manifest symmetry between the left- and right-handed sectors, the right-handed quark mixing is not necessarily the same as the left-handed quark mixing described by the CKM matrix. Thus we have additional right-handed charged current interactions with couplings different from the left ones, which are suppressed by the heavy mass of the extra W boson. The strength of the right-handed couplings should be determined by measurements in various phenomena. On the other hand, when the electroweak symmetry is dynamically broken, some non-universal interactions may exist which lead to additional right-handed and left-handed couplings on charged current interactions [8]. If the anomalous right-handed $\bar{t}bW$ couplings exist, their effects can be found in rare B decays [9,10] and also in the various phenomena at future colliders [11,12].

The production of 10^7–10^8 top quark pairs per year expected at the Large Hadron Collider (LHC) will allow us to study the structure of the top quark couplings. The $\bar{t}bW$ coupling will be directly measured with high precision through the dominant $t \to bW$ channel and the anomalous $\bar{t}bW$ coupling will be tested in a direct way. The subdominant channel of the top quark is the CKM non-diagonal decay $t \to sW$ in the SM, of which the branching ratio is estimated to be $\mathrm{Br}(t \to sW) \sim 1.6 \times 10^{-3}$, when $|V_{ts}| = 0.04$ is assumed. Despite the small branching ratio of this channel, the large number of expected top quark productions at LHC will enable us to measure the $t \to sW$ process and provide us with a chance to probe the $\bar{t}sW$ coupling directly. Hence the anomalous $\bar{t}sW$ coupling is worth studying at present. In this work, we consider the

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anomalous right-handed $\bar{t}bW$ and $\bar{t}sW$ couplings and their impact on the CP violation in $B \to \phi K_S$ decay.

We do not specify the underlying model here but concentrate on the anomalous right-handed couplings of charged current interactions, ignoring the effects of additional left-handed interactions and new particles. The relevant right-handed couplings are described by the effective Lagrangian

$$\mathcal{L} = -\frac{g}{\sqrt{2}} \sum_{q=s,b} V_{tq}\ \bar{t}\gamma^\mu (P_\mathrm{L} + \xi_q P_\mathrm{R}) q W_\mu^+ + \mathrm{H.c.}, \quad (3)$$

where ξ_q is a dimensionless parameter measuring new physics effects. If ξ_q has a complex phase, generically it invokes a new CP violation leading to a shift of the observed $\sin 2\beta$.

This paper is organized as follows. In Sect. 2, the effective Hamiltonian formalism with right-handed $\bar{t}bW$ and $\bar{t}sW$ couplings is given. In Sect. 3, we discuss the constraints on the parameters ξ_b and ξ_s using the radiative $B \to X_s\gamma$ decay. The analysis on hadronic decays $B \to J/\psi K$ and $B \to \phi K$ is presented in Sect. 4 to extract the corresponding $\sin 2\beta$. Finally we conclude in Sect. 5.

2 The effective Hamiltonian

The effective Hamiltonian approach is required when we study rare decays of B mesons in order to incorporate QCD effects in a consistent way. The $\Delta B = 1$ effective Hamiltonian for describing hadronic B decays is given by

$$\begin{aligned}\mathcal{H}_\mathrm{eff} = \frac{4G_\mathrm{F}}{\sqrt{2}} V_{ts}^* V_{tb} \Bigg[&\sum_{i=1}^{2} (C_i(\mu) O_i(\mu) + C_i'(\mu) O_i'(\mu)) \\ &- \sum_{i=3}^{10} (C_i(\mu) O_i(\mu) + C_i'(\mu) O_i'(\mu)) \Bigg] \\ &+ \mathrm{H.c.}, \end{aligned} \quad (4)$$

including the effects of the anomalous right-handed top quark interactions. The operator basis is defined following [13] by

$$\begin{aligned}
O_1 &= (\bar{s}_\alpha c_\beta)_\mathrm{L} (\bar{c}_\beta b_\alpha)_\mathrm{L}, \\
O_2 &= (\bar{s}c)_\mathrm{L} (\bar{c}b)_\mathrm{L}, \\
O_3 &= (\bar{s}b)_\mathrm{L} \sum_{q'=u,d,s,c,b} (\bar{q}'q')_\mathrm{L}, \\
O_4 &= (\bar{s}_\alpha b_\beta)_\mathrm{L} \sum_{q'=u,d,s,c,b} (\bar{q}'_\beta q'_\alpha)_\mathrm{L}, \\
O_5 &= (\bar{s}b)_\mathrm{L} \sum_{q'=u,d,s,c,b} (\bar{q}'q')_\mathrm{R}, \\
O_6 &= (\bar{s}_\alpha b_\beta)_\mathrm{L} \sum_{q'=u,d,s,c,b} (\bar{q}'_\beta q'_\alpha)_\mathrm{R}, \\
O_7 &= \frac{3}{2} (\bar{s}b)_\mathrm{L} \sum_{q'=u,d,s,c,b} e_{q'} (\bar{q}'q')_\mathrm{L}, \\
O_8 &= \frac{3}{2} (\bar{s}_\alpha b_\beta)_\mathrm{L} \sum_{q'=u,d,s,c,b} e_{q'} (\bar{q}'_\beta q'_\alpha)_\mathrm{L}, \\
O_9 &= \frac{3}{2} (\bar{s}b)_\mathrm{L} \sum_{q'=u,d,s,c,b} e_{q'} (\bar{q}'q')_\mathrm{R}, \\
O_{10} &= \frac{3}{2} (\bar{s}_\alpha b_\beta)_\mathrm{L} \sum_{q'=u,d,s,c,b} e_{q'} (\bar{q}'_\beta q'_\alpha)_\mathrm{R}, \\
O_{11} &= \frac{g_s}{16\pi^2} m_b \bar{s}_\alpha P_\mathrm{R} \sigma_{\mu\nu} T^a_{\alpha\beta} b_\beta G^{a\ \mu\nu}, \\
O_{12} &= \frac{e}{16\pi^2} m_b \bar{s} P_\mathrm{R} \sigma_{\mu\nu} b F^{\mu\nu},
\end{aligned} \quad (5)$$

where $(\bar{q}b)_\mathrm{L/R} = (\bar{q}\gamma_\mu P_\mathrm{L/R} b)$. The operators O_i' are the chiral conjugates of the O_i operators.

Matching the effective Hamiltonian and our model Lagrangian of (2) at the $\mu = m_W$ scale, we have the Wilson coefficients $C_i(\mu = m_W)$ and $C_i'(\mu = m_W)$ in the SM:

$$\begin{aligned}
C_1(m_W) &= \frac{11}{2} \frac{\alpha_\mathrm{s}(m_W)}{4\pi}, \\
C_2(m_W) &= 1 - \frac{11}{6} \frac{\alpha_\mathrm{s}(m_W)}{4\pi} - \frac{35}{18} \frac{\alpha}{4\pi}, \\
C_3(m_W) &= -\frac{\alpha_\mathrm{s}(m_W)}{24\pi} E_0(x_t) \\
&\quad + \frac{\alpha}{6\pi} \frac{1}{\sin^2\theta_\mathrm{W}} [2\ B_0(x_t) + C_0(x_t)], \\
C_4(m_W) &= \frac{\alpha_\mathrm{s}(m_W)}{8\pi} E_0(x_t), \\
C_5(m_W) &= -\frac{\alpha_\mathrm{s}(m_W)}{24\pi} E_0(x_t), \\
C_6(m_W) &= \frac{\alpha_\mathrm{s}(m_W)}{8\pi} E_0(x_t), \\
C_7(m_W) &= \frac{\alpha}{6\pi} [4\ C_0(x_t) + D_0(x_t)], \\
C_9(m_W) &= \frac{\alpha}{6\pi} \Bigg[4\ C_0(x_t) + D_0(x_t) \\
&\quad + \frac{1}{\sin^2\theta_\mathrm{W}} (10\ B_0(x_t) - 4\ C_0(x_t)) \Bigg], \\
C_8(m_W) &= C_{10}(m_W) = 0, \\
C_{11}(m_W) &= G(x_t), \\
C_{12}(m_W) &= F(x_t), \\
C_i'(m_W) &= 0, \qquad i = 1, \cdots, 12,
\end{aligned} \quad (6)$$

where $B_0(x), C_0(x), D_0(x), E_0(x), F(x)$, and $G(x)$ are the well-known Inami–Lim loop functions of which the explicit forms are given in [13,14]. Turning on the right-handed $\bar{t}bW$ and $\bar{t}sW$ couplings, we have the modification of the loop functions in the Wilson coefficients C_i,

$$\begin{aligned}
D_0(x_t) &\to D_0(x_t) + \xi_b \frac{m_b}{m_t} D_\mathrm{R}(x_t), \\
E_0(x_t) &\to E_0(x_t) + \xi_b \frac{m_b}{m_t} E_\mathrm{R}(x_t), \\
F(x_t) &\to F(x_t) + \xi_b \frac{m_t}{m_b} F_\mathrm{R}(x_t), \\
G(x_t) &\to G(x_t) + \xi_b \frac{m_t}{m_b} G_\mathrm{R}(x_t),
\end{aligned} \quad (7)$$

and we also have the new Wilson coefficients C_i':

$$
\begin{aligned}
C_3'(m_W) &= -\frac{\alpha_s(m_W)}{24\pi}\xi_s\frac{m_b}{m_t}E_{\mathrm{R}}(x_t),\\
C_4'(m_W) &= \frac{\alpha_s(m_W)}{8\pi}\xi_s\frac{m_b}{m_t}E_{\mathrm{R}}(x_t),\\
C_5'(m_W) &= -\frac{\alpha_s(m_W)}{24\pi}\xi_s\frac{m_b}{m_t}E_{\mathrm{R}}(x_t),\\
C_6'(m_W) &= \frac{\alpha_s(m_W)}{8\pi}\xi_s\frac{m_b}{m_t}E_{\mathrm{R}}(x_t),\\
C_7'(m_W) &= \frac{\alpha}{6\pi}\xi_s\frac{m_b}{m_t}D_{\mathrm{R}}(x_t),\\
C_9'(m_W) &= \frac{\alpha}{6\pi}\xi_s\frac{m_b}{m_t}D_{\mathrm{R}}(x_t),\\
C_{11}'(m_W) &= \xi_s\frac{m_t}{m_b}G_{\mathrm{R}}(x_t),\\
C_{12}'(m_W) &= \xi_s\frac{m_t}{m_b}F_{\mathrm{R}}(x_t),\\
C_i'(m_W) &= 0, \quad (i=1,2,8,10),
\end{aligned} \tag{8}
$$

where the new loop functions are given by

$$
\begin{aligned}
D_{\mathrm{R}}(x) &= \frac{x(59-38x+25x^2+2x^3)}{36(x-1)^4}+\frac{2(x+1)}{3(x-1)^5}\ln x\\
&\quad+\frac{x^2}{2(x-1)^4}\ln x,\\
E_{\mathrm{R}}(x) &= \frac{x(-116+165x-114x^2+29x^3)}{18(x-1)^4}\\
&\quad+\frac{2+3x+x^2}{3(x-1)^5}\ln x,\\
F_{\mathrm{R}}(x) &= \frac{-20+31x-5x^2}{12(x-1)^2}+\frac{x(2-3x)}{2(x-1)^3}\ln x,\\
G_{\mathrm{R}}(x) &= -\frac{4+x+x^2}{4(x-1)^2}+\frac{3x}{2(x-1)^3}\ln x,
\end{aligned} \tag{9}
$$

where our new loop functions $F_{\mathrm{R}}(x)$, $G_{\mathrm{R}}(x)$ and $D_{\mathrm{R}}(x)$ agree with those in [15,16] and $E_{\mathrm{R}}(x)$ is the first calculation. Note that the $\mathcal{O}(\xi)$ terms of the Z penguin diagram are suppressed by the heavy mass of the Z boson as m_b^2/m_Z^2 or q^2/m_Z^2, and we neglect them here. For the box diagram, if we include only one anomalous coupling, the chirality structures of the two currents are different and one current is proportional to the fermion momentum and the other current is proportional to the fermion mass. This indicates that the integrand is always an odd function and the loop integral vanishes. Therefore the $\mathcal{O}(\xi)$ terms of the box diagram do not exist and the leading contribution is of order ξ^2. Hence, we also ignore the box contribution. As a consequence, the contribution of order $\mathcal{O}(\xi)$ comes only through the γ penguin and gluon penguin diagrams. Actually the contributions of the $O_{12}^{(\prime)}$ operator to hadronic decays are very small and we neglect them in the numerical analysis.

The renormalization group (RG) evolution of the Wilson coefficients $\mathbf{C}=(C_i, C_i')^\dagger$ given by

$$
\mu\frac{\mathrm{d}}{\mathrm{d}\mu}\mathbf{C}(M_W) = -\frac{g_s^2}{16\pi^2}\gamma^{\mathrm{T}}\mathbf{C}(M_W)
$$

is governed by a 24×24 anomalous dimension matrix γ. Since the strong interaction preserves chirality, the operators O_i and O_i' are evolved separately without mixing between them. Thus the 24×24 anomalous dimension matrix γ is decomposed into two identical 12×12 matrices γ_0 given in the SM. The 12×12 anomalous dimension matrix γ_0 can be found in [13,17,18]. The evolved Wilson coefficients $C_i^{(\prime)}(\mu)$ are expressed in terms of the initial conditions of (5) and (7), $C(\mu)=U(\mu,M_W)C(M_W)$. The explicit formula for the evolution matrix $U(\mu,M_W)$ can be found in [17,18]. The matrix elements of the operators also have one loop corrections. We define the effective Wilson coefficients by absorbing the correction of the matrix elements in the Wilson coefficients as given in [19–21]. Then the Hamiltonian is expressed in terms of effective Wilson coefficients and tree level matrix elements.

3 $B \to X_s\gamma$ constraints

Before the analysis of $\sin 2\beta$, we consider the radiative $B\to X_s\gamma$ decay to constrain the model. This channel has already been observed experimentally and a more precise measurement will be obtained from the accumulation of the data at B factories. It is well known that this process is an effective probe of new physics since the dominant penguin diagram is sensitive to the internal heavy particle property. Especially, the right-handed couplings inside the loop of the operators O_{11} and O_{12} involve an enhancement factor m_t/m_b. Thus stringent limits on ξ_b and ξ_s are yielded from the measurement of the $B\to X_s\gamma$ decay [15, 9]. We present the updated constraints on the anomalous couplings from the branching ratio and the bound of CP violating asymmetry in $B\to X_s\gamma$ decay.

The weighted average of the branching ratio is given by

$$
\mathrm{Br}(B\to X_s\gamma) = (3.23\pm 0.41)\times 10^{-4}, \tag{10}
$$

from the measurements of Belle [22], CLEO [23] and ALEPH [24] groups. The CP violating asymmetry in the $B\to X_s\gamma$ decays defined by

$$
A_{CP}(B\to X_s\gamma) = \frac{\Gamma(\bar{B}\to X_s\gamma)-\Gamma(B\to X_{\bar{s}}\gamma)}{\Gamma(\bar{B}\to X_s\gamma)+\Gamma(B\to X_{\bar{s}}\gamma)}
$$

is very small in the SM because of the unitarity of the CKM matrix. The direct CP asymmetry A_{CP} is measured by CLEO [25]:

$$
\begin{aligned}
&A_{CP}(B\to X_s\gamma)\\
&= (-0.079\pm 0.108\pm 0.022)(1.0\pm 0.030),
\end{aligned} \tag{11}
$$

where the first error is statistical, the second one is additive systematic over the various $b\to s\gamma$ decay modes, and the third one is multiplicative systematic. Note that the present measurement of A_{CP} is still consistent with 0. The complex anomalous $\bar{t}bW$ coupling contributes to the CP asymmetry through the interference terms of the O_{11} and O_{12} operators such as $\delta A_{CP}\sim a_1\mathrm{Im}C_2C_{12}^* + a_2\mathrm{Im}C_{11}C_{12}^* +$

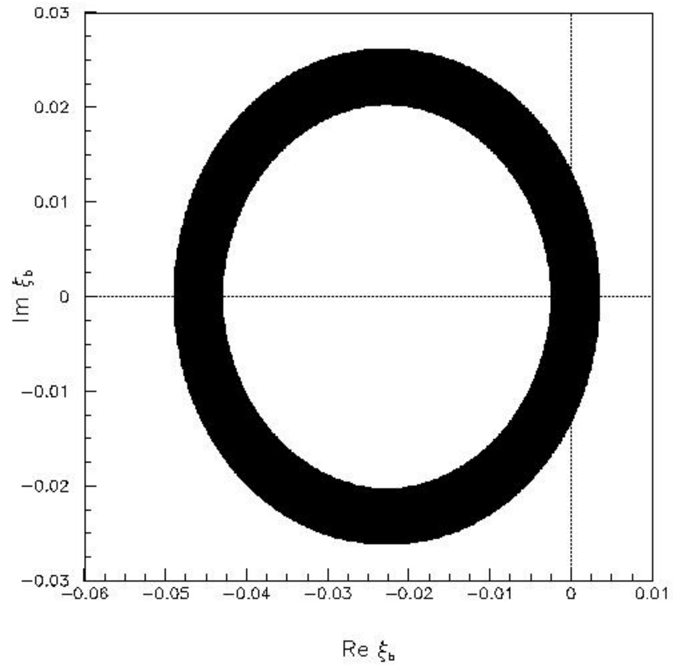


Fig. 1. Allowed parameter set ($\mathrm{Re}\xi_b, \mathrm{Im}\xi_b$) under the constraints by the branching ratio and CP asymmetry in $B \to X_s\gamma$ decay

$a_3 \mathrm{Im} C_2 C_{11}^*$, which provides an additional test on ξ_b, independent of the branching ratio. On the contrary, the $\bar{t}sW$ coupling does not contribute to A_{CP} at this level.

The explicit expressions of the branching ratio and the CP asymmetry are presented in [26,27] in terms of the evolved Wilson coefficients at the $\mu = m_b$ scale. With the measured values of (9) and (10), we obtain the constraints on ξ_b at 2σ C.L. as

$$-0.002 < \mathrm{Re}\xi_b + 22|\xi_b|^2 < 0.0033, \tag{12}$$
$$-0.299 < \frac{0.27\ \mathrm{Im}\xi_b}{0.095 + 12.54\mathrm{Re}\xi_b + 414.23|\xi_b|^2} < 0.141,$$

and the allowed parameter set ($\mathrm{Re}\xi_b, \mathrm{Im}\xi_b$) is depicted in Fig. 1. Since ξ_s is irrelevant for the CP asymmetry, we can set the limit on ξ_s to be

$$|\xi_s| < 0.012, \tag{13}$$

from the branching ratio alone [15].

4 Hadronic decays

4.1 $B \to J/\psi K$

The $B \to J/\psi K$ decays are dominated by the tree level $b \to c\bar{c}s$ decay amplitude and a single weak phase in the SM. The subleading penguin contribution depends on the CKM factor $V_{tb}V_{ts}^*$ which gives the same phase as the factor $V_{cb}V_{cs}^*$ of the tree diagram and the weak phase structure is not affected. On that account, this mode is thought to be a golden mode to extract the weak phase β. The CP asymmetries in $B \to J/\psi K$ decays given in (2), $\sin 2\beta_{\mathrm{eff}} = 0.734 \pm 0.054$, agree well with the SM prediction. The subscript "eff" denotes the "observed" $\sin 2\beta$. In terms of the Wilson coefficients, the decay amplitude for $B \to J/\psi K_S$ decay is dominated by C_2 which involves no $\xi_{b,s}$ effects. The subdominant amplitude involving $\xi_{b,s}$ is suppressed by loop suppression and/or a CKM factor as well as the $\xi_{b,s}$ itself, of which the suppression factor is estimated to be of order $< 10^{-4}$. Thus a new physics effect on the decay amplitude is ignored to a very good approximation. Considering the B–$\bar{B}$ mixing with right-handed coupling, $\mathcal{O}(\xi_b)$ contributions vanish in the box diagram calculation by the chirality relation and the leading new physics contribution to the off-diagonal matrix element is of order $\mathcal{O}(\xi_b^2)$,

$$M_{12} = M_{12}^{\mathrm{SM}} \left(1 + \xi_b^2 \frac{S_{\mathrm{R}}(x_t)}{S_0(x_t)} \frac{(\bar{b}P_{\mathrm{L}}d)(\bar{b}P_{\mathrm{L}}d)}{(\bar{b}\gamma_\mu P_{\mathrm{L}}d)(\bar{b}\gamma^\mu P_{\mathrm{L}}d)}\right), \tag{14}$$

where the new loop function $S_{\mathrm{R}}(x)$ is given by

$$S_{\mathrm{R}}(x) = \frac{x(x^2 - 2x + 6)}{(1-x)^2} + \frac{x(x+2)(x^2 - x + 2)}{(1-x)^3} \ln x, \tag{15}$$

and the SM loop function $S_0(x)$ can be found in [13]. This leads to the $\mathcal{O}(\xi_b^2)$ shift of the weak phase

$$\sin 2\beta_{\mathrm{eff}} = \sin 2\beta + 4.3|\xi_b|^2 \sin 2\varphi, \tag{16}$$

where $\xi_b = |\xi_b| \mathrm{e}^{\mathrm{i}\varphi}$ and

$$\frac{\langle B^0|(\bar{b}P_{\mathrm{L}}d)(\bar{b}P_{\mathrm{L}}d)|\bar{B}^0\rangle}{\langle B^0|(\bar{b}\gamma_\mu P_{\mathrm{L}}d)(\bar{b}\gamma^\mu P_{\mathrm{L}}d)|\bar{B}^0\rangle} \approx \frac{3}{4}\left(\frac{m_B}{m_b}\right)^2.$$

With the allowed parameter set of Fig. 1, the second term of (16) is at most of order 10^{-3} so we can neglect it for the discussion of $\sin 2\beta$. On the other hand, the anomalous $\bar{t}sW$ coupling is irrelevant for the B–$\bar{B}$ mixing and invokes no effects on $\sin 2\beta_{\mathrm{eff}}$ in $B \to J/\psi K$ decays. As a consequence, the observed weak phase $\sin 2\beta$ in $B \to J/\psi K$ decays is hardly affected by the right-handed top couplings.

4.2 $B \to \phi K_s$

The average of the CP asymmetry in the $B \to \phi K_S$ decay measured by the BaBar [1] and Belle [2] groups is given by

$$A_{CP}^{\phi K} = -0.56 \pm 0.43,$$
$$\sin 2\beta_{\mathrm{eff}}^{\phi K} = -0.39 \pm 0.41, \tag{17}$$

where $A_{CP}^{\phi K}$ is the CP violating asymmetry defined by $A_{CP}^{\phi K} \equiv [\Gamma(B \to \phi K) - \Gamma(\bar{B} \to \phi\bar{K})]/[\Gamma(B \to \phi K) + \Gamma(\bar{B} \to \phi\bar{K})]$ and $\sin 2\beta_{\mathrm{eff}}^{\phi K}$ the observed weak phase extracted from $B \to \phi K_S$ decay.

The $b \to s\bar{s}s$ transition responsible for the $B \to \phi K$ decays arises at one loop level in the SM. It is known that the gluon penguin diagram plays a central role in this decay channel through the chromo-magnetic (dipole penguin) operator O_{11} as well as the four quark operator. As in the case of $B \to X_s\gamma$ decay, if the right-handed couplings are switched on, the enhancement factor m_t/m_b involved in the penguin loop makes the new effect of the $O_{11}^{(\prime)}$ operator lead to significant contributions in $B \to \phi K$ decays. It has been discussed that the anomalous right-handed $\bar{t}bW$ coupling can yield a deviation of the CP asymmetry in the $B \to \phi K_S$ process from the SM prediction by Abd El-Hady and Valencia [10]. Here we present the detailed analysis on the CP violation in $B \to \phi K_S$ decays including both of the $\bar{t}bW$ and $\bar{t}sW$ couplings and compare $A_{CP}^{\phi K}$ and $\sin 2\beta_{\mathrm{eff}}^{\phi K}$ with the experiment. On the other hand, the electroweak penguin operators also give a sizable contribution to this decay mode: up to 20% [19]. Therefore we include all operators in the effective Hamiltonian to evaluate the $B \to \phi K_S$ decay rate except for O_{12} since its contribution is extremely small.

With the definition of the form factors and decay constants

$$
\begin{aligned}
\langle P(p')|V_\mu|B(p)\rangle &= \left[(p+p')_\mu - \frac{m_B^2-m_P^2}{q^2}q_\mu\right]F_1^P(q^2) \\
&\quad + \frac{m_B^2-m_P^2}{q^2}q_\mu F_0^P(q^2), \\
\langle 0|A_\mu|P(p)\rangle &= \mathrm{i} f_P p_\mu, \\
\langle 0|V_\mu|V(p)\rangle &= f_V m_V \epsilon_\mu,
\end{aligned} \tag{18}
$$

we write the decay amplitude for $B \to \phi K$ decays as

$$
\begin{aligned}
\mathcal{A}(B^0 \to \phi K^0) &= -\frac{G_\mathrm{F}}{\sqrt{2}}V_{tb}^* V_{ts} \\
&\quad \times \left[a_3+a_4+a_5-\frac{1}{2}(a_7+a_9+a_{10})\right] \\
&\quad \times 2 f_\phi m_\phi (\epsilon^* \cdot p_B) F_1^K + A_{11}^{\phi K},
\end{aligned} \tag{19}
$$

where $a_{2i-1} = C_{2i-1} + C_{2i}/N_c$, $a_{2i} = C_{2i} + C_{2i-1}/N_c$. The contribution of the chromo-magnetic operator $A_{11}^{\phi K}$ is given by [28,29]

$$
\begin{aligned}
A_{11}^{\phi K} &\equiv \langle \phi K^0|(\mathcal{H}_{11}+\mathcal{H'}_{11})|B^0\rangle, \\
&= \frac{G_\mathrm{F}}{\sqrt{2}}\frac{\alpha_\mathrm{s}(q^2)}{4\pi q^2}V_{tb}^* V_{ts} m_b(\mu)\frac{N_c^2-1}{N_c^2} f_\phi m_\phi (\epsilon^*\cdot p_B) \\
&\quad \times (C_{11}+C'_{11})(F_1^K X + F_0^K Y),
\end{aligned} \tag{20}
$$

with

$$
\begin{aligned}
X &= 4m_b + 5m_s + 3m_s\left(\frac{m_B^2-m_K^2}{m_\phi^2}\right) \\
&\quad - \left(\frac{3m_B^2-3m_K^2+m_\phi^2}{8m_b}\right)\left(1+\frac{m_B^2-m_K^2}{m_\phi^2}\right), \\
Y &= \frac{3}{2}\left(\frac{m_B^2-m_K^2}{m_b-m_s}\right) \\
&\quad + \left(\frac{m_B^2-m_K^2}{m_\phi^2}\right)\left(\frac{3m_B^2-3m_K^2+m_\phi^2}{8m_b} - 3m_s\right),
\end{aligned} \tag{21}
$$

and $q^2 = m_B^2/2 - m_K^2/4 + m_\phi^2/2$. The B to K form factor, $F_{0,1}$, is the principal source of hadronic uncertainty for this process. An early calculation was performed in the framework of the quark model [30]. We can set $F_0 = F_1$ close to the point $q^2 = 0$ [31] and assume simple pole dominance. Here we take the value of $F_{0,1}(0) = 0.26$–0.37 from the QCD sum rule results [32]. Note that new effects on the four quark operators are doubly suppressed by both m_b/m_t and $\xi_{b,s}$, while the effects on the dipole operator involve an enhancement factor m_t/m_b compensating the $\xi_{b,s}$ suppression. Thus the new contribution dominantly comes through $A_{11}^{\phi K}$. We also note that $C_i^{(\prime)}$ in (19) and (20) are the effective Wilson coefficients absorbing the one loop correction to the hadronic matrix elements and they involve the strong phases.

Since the four quark operator contribution in the first term in (19) involves a strong phase, the new phase of $A_{11}^{\phi K}$ leads to a deviation of $|\bar{A}/A|$ from unity and we have the rate asymmetry implying the direct CP violation. In terms of the parameter λ defined by

$$
\lambda = \sqrt{\frac{M_{12}^*}{M_{12}}}\frac{\bar{A}}{A}, \tag{22}
$$

where $A = \mathcal{A}(B^0 \to \phi K^0)$ and $\bar{A} = \mathcal{A}(\bar{B}^0 \to \phi \bar{K}^0)$, we write the full expression of the time-dependent CP asymmetry as

$$
\begin{aligned}
a_{\phi K}(t) &\equiv \frac{\Gamma(B_\mathrm{phys}^0(t)\to\phi K^0) - \Gamma(\bar{B}_\mathrm{phys}^0(t)\to\phi\bar{K}^0)}{\Gamma(B_\mathrm{phys}^0(t)\to\phi K^0) + \Gamma(\bar{B}_\mathrm{phys}^0(t)\to\phi\bar{K}^0)}, \\
&= C_{\phi K}\cos\Delta m_B t + S_{\phi K}\sin\Delta m_B t,
\end{aligned} \tag{23}
$$

where the coefficients are

$$
\begin{aligned}
C_{\phi K} &= \frac{1-|\lambda|^2}{1+|\lambda|^2} \equiv -A_{CP}^{\phi K}, \\
S_{\phi K} &= -\frac{2\mathrm{Im}\lambda}{1+|\lambda|^2} \equiv \sin 2\beta_\mathrm{eff}^{\phi K}.
\end{aligned} \tag{24}
$$

Note that the hadronic uncertainty is cancelled in λ and the CP violating observables $A_{CP}^{\phi K}$ and $\sin 2\beta_\mathrm{eff}^{\phi K}$ are free from the hadronic uncertainty. We can express the parameter λ by

$$
\begin{aligned}
\lambda &= \lambda^\mathrm{SM}\left(\frac{1+21.84\ \mathrm{e}^{-\mathrm{i}\delta}\ \xi_q}{1+21.84\ \mathrm{e}^{-\mathrm{i}\delta}\ \xi_q^*}\right) \\
&\approx \lambda^\mathrm{SM}(1+\mathrm{i}\ 43.7\ |\xi_q|\ \mathrm{e}^{-\mathrm{i}\delta}\sin\varphi_q),
\end{aligned} \tag{25}
$$

where φ_q is the phase of ξ_q, and δ the strong phase introduced by the one loop correction to the matrix elements. For $\lambda^\mathrm{SM} \equiv \mathrm{e}^{\mathrm{i}\beta_\mathrm{SM}}$, we will use the measured value given in (2). With this expression, we can write the CP asymmetries as

$$
\begin{aligned}
A_{CP}^{\phi K} &= 23.3\ |\xi_q|\ \sin\varphi_q, \\
\sin 2\beta_\mathrm{eff}^{\phi K} &= \sin 2\beta + 52.2\ |\xi_q|\ \sin\varphi_q,
\end{aligned} \tag{26}
$$

where $\delta = 2.58$ in our calculation. Note that the second expression of (25) is no more valid for the maximal value

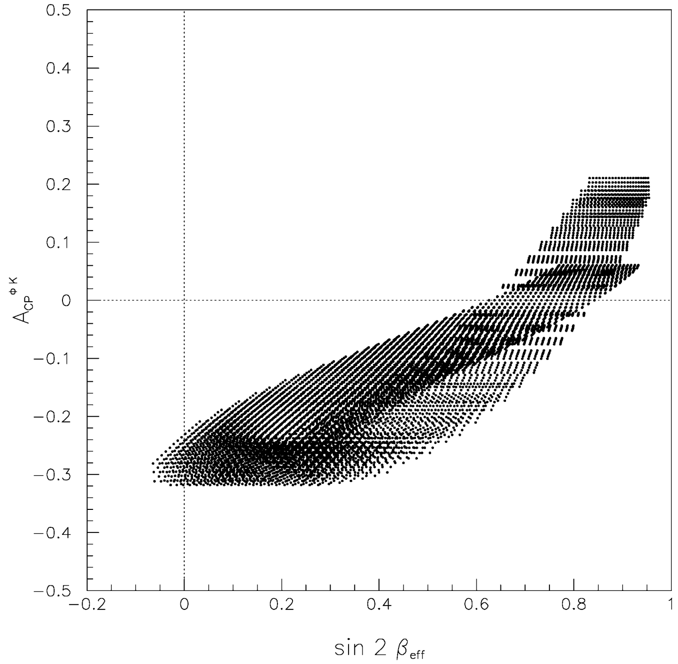


Fig. 2. Correlation of $\sin 2\beta_{\rm eff}$ and $A_{CP}^{\phi K}$ with varying ξ_b as shown in Fig. 1

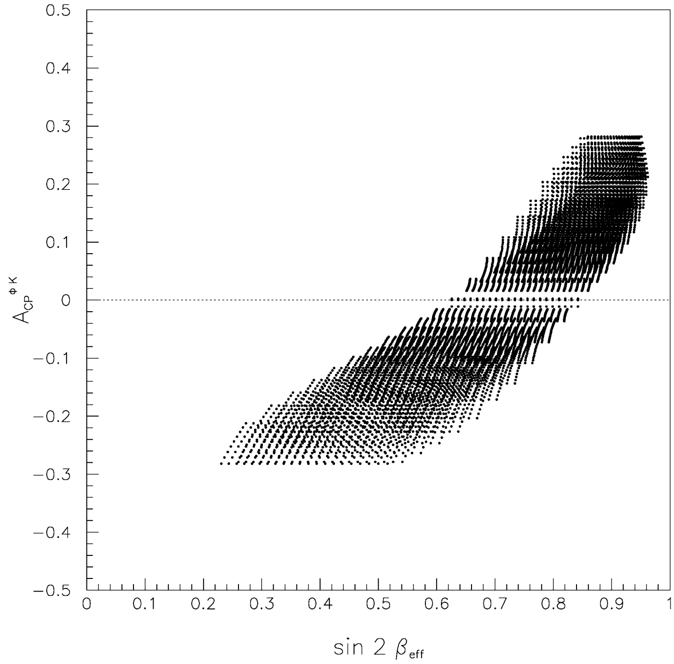


Fig. 3. Correlation of $\sin 2\beta_{\rm eff}$ and $A_{CP}^{\phi K}$ with varying ξ_s under the constraint of (11)

of $|\xi_b| \sim 0.04$ and so are the above expressions of $A_{CP}^{\phi K}$ and $\sin 2\beta_{\rm eff}^{\phi K}$.

With the allowed parameter set given in Fig. 1 and the $\sin 2\beta$ measured in $B \to J/\psi K$ decay, given in (13), we have the rate asymmetry $A_{CP}^{\phi K}$ and the effective weak phase $\sin 2\beta_{\rm eff}^{\phi K}$

$$\begin{aligned} &-0.34 < A_{CP}^{\phi K} < 0.22, \\ &-0.10 < \sin 2\beta_{\rm eff}^{\phi K} < 0.96. \end{aligned} \tag{27}$$

As shown in the previous section, the effect of the right-handed top couplings on the B–$\bar{B}$ mixing sector is safely neglected for the evaluation of $\sin 2\beta_{\rm eff}$. In Fig. 2, the correlation of $\sin 2\beta_{\rm eff}^{\phi K}$ and $A_{CP}^{\phi K}$ is shown. We find that a large rate asymmetry ($-20 \sim -30\%$) should exist for the observed $\sin 2\beta$ to be negative. Even if the future experiments ascertain that the $\sin 2\beta_{\rm eff}^{\phi K}$ is consistent with the SM prediction, it is still possible that there exists a sizable direct CP violation $A_{CP}^{\phi K} \sim 10\%$. With the right-handed $\bar{t}sW$ coupling $|\xi_s| < 0.012$, we have

$$\begin{aligned} &-0.28 < A_{CP}^{\phi K} < 0.28, \\ &\ \ 0.23 < \sin 2\beta_{\rm eff}^{\phi K} < 0.94, \end{aligned} \tag{28}$$

and their correlation is shown in Fig. 3. We also find that the large CP asymmetry ($\sim 10\%$) is possible with ξ_s even if $\sin 2\beta_{\rm eff}$ agrees with the SM prediction.

We also calculate the branching ratio:

$$\mathrm{Br}(B \to \phi K) = \tau_B \frac{1}{16\pi} \frac{\lambda(m_B^2, m_\phi^2, m_K^2)}{m_B^3} |\mathcal{A}|^2, \tag{29}$$

where $\lambda(x,y,z) = (x^2+y^2+z^2-2xy-2yz-2zx)^{1/2}$ and τ_B is the lifetime of the B meson. Figures 4 and 5 show the relations of the branching ratio and CP violations in the presence of the right-handed $\bar{t}bW$ and $\bar{t}sW$ couplings. The present measurements of the branching ratio for the $B^0 \to \phi K^0$ decay read

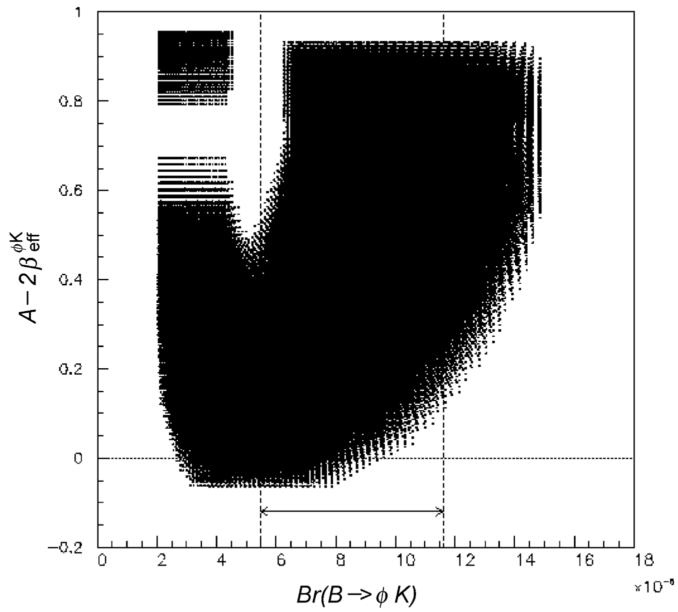
$$\begin{aligned} &\mathrm{Br}(B^0 \to \phi K^0) \\ &= (5.4^{+3.7}_{-2.7} \pm 0.7) \times 10^{-6} < 12.3 \times 10^{-6} \quad \text{CLEO}, \\ &= (8.1^{+3.1}_{-2.5} \pm 0.8) \times 10^{-6} \quad \text{BaBar}, \\ &= (8.7^{+3.8}_{-3.0} \pm 1.5) \times 10^{-6} \quad \text{Belle}, \end{aligned} \tag{30}$$

given by the CLEO [35], BaBar [36] and Belle [37] groups. Since the CLEO result is just an intermediate fitted value and the Belle result is a preliminary one, we just show the BaBar result in Figs. 4 and 5. In our evaluation, the SM value is (1.9–4.0) $\times 10^{-6}$, close to the prediction of [38]. We can see that the negative $\sin 2\beta_{\rm eff}$ consistent with the BaBar cross section is possible with the anomalous $\bar{t}bW$ coupling, but we do not expect such a solution with the anomalous $\bar{t}sW$ coupling.

We also show the correlation of the CP asymmetries between the $B \to X_s\gamma$ and $B \to \phi K$ decays in Fig. 6. For the negative $\sin 2\beta_{\rm eff}$ in $B \to \phi K$ decay, $-(2$–$3)\%$ of A_{CP}^{γ} is expected.

5 Concluding remarks

We have studied the effects of the complex right-handed top quark couplings on the CP violation in $B \to \phi K$ decays, which originate in the general $SU(2)_{\rm L} \times SU(2)_{\rm R} \times$



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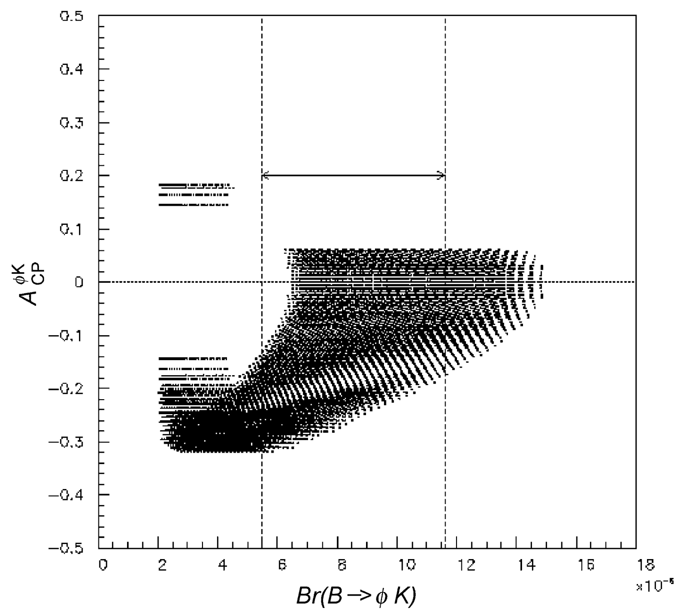
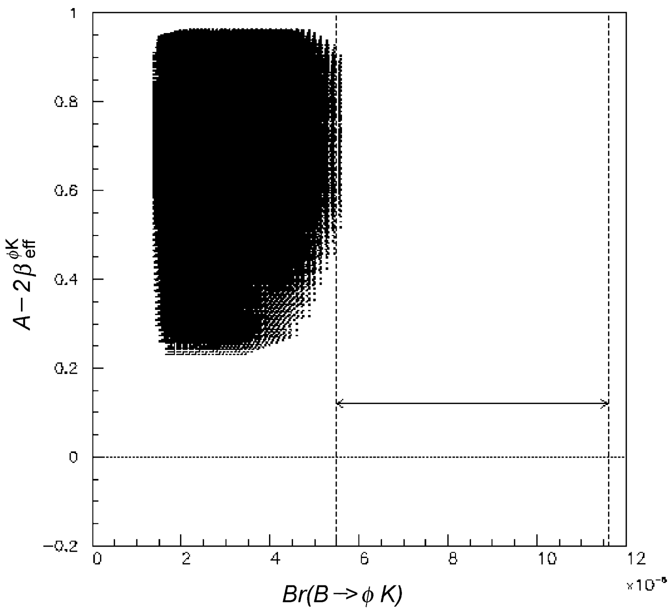


b

Fig. 4. **a** Correlation of the branching ratio of $B \to \phi K$ decay and $\sin 2\beta^{\phi K}_{\rm eff}$ with varying ξ_b as shown in Fig. 1. **b** Correlation of the branching ratio of $B \to \phi K$ decay and $A^{\phi K}_{CP}$ with varying ξ_b as shown in Fig. 1



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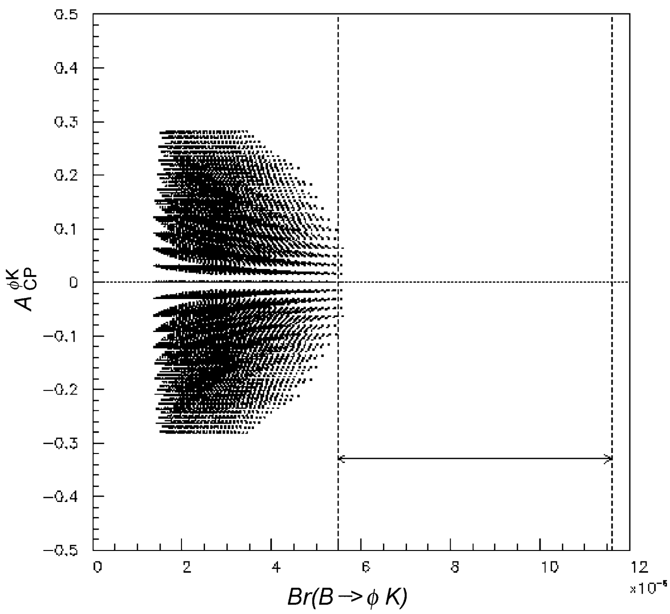
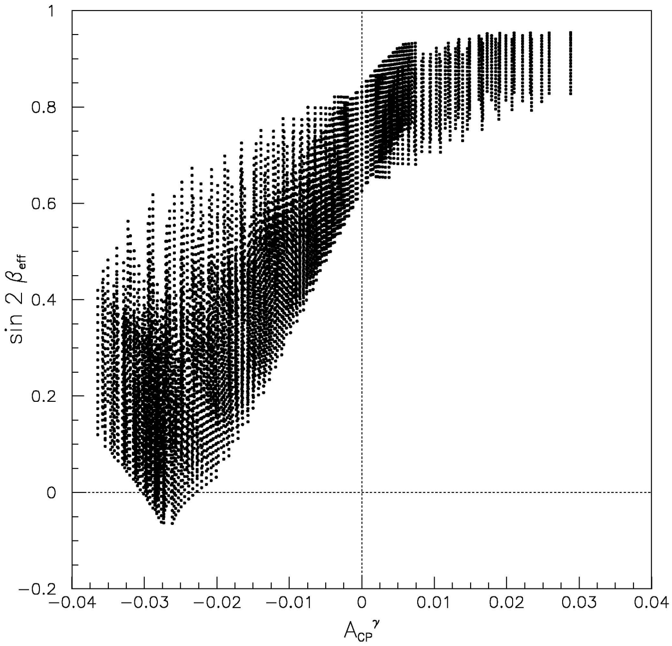


b

Fig. 5. **a** Correlation of the branching ratio of $B \to \phi K$ decay and $\sin 2\beta^{\phi K}_{\rm eff}$ with varying ξ_s under the constraint of (11). **b** Correlation of the branching ratio of $B \to \phi K$ decay and $A^{\phi K}_{CP}$ with varying ξ_s under the constraint of (11)

$U(1)$ model or the dynamical electroweak symmetry breaking model. Since the contribution of those couplings to the B–$\bar{B}$ mixing is suppressed by the quadratic order of ξ_q, the measurement of the $\sin 2\beta$ in $B \to J/\psi K$ decays is not affected by the right-handed couplings. However, the gluonic dipole penguin operator, which plays a important role in $b \to s\bar{s}s$ decay, gets a sizable contribution from the right-handed couplings due to an enhancement factor m_t/m_b and the observed $\sin 2\beta$ in $B \to \phi K$ decays can be shifted. Even a negative $\sin 2\beta$ is possible with the anomalous $\bar{t}bW$ coupling, as is consistent with the recent measurements. Note that the value of $\sin 2\beta$ with the anomalous $\bar{t}sW$ coupling is also shifted but still positive due to a stricter bound on $|\xi_s|$ than that on $|\xi_b|$. In conclusion, the right-handed top couplings are good candidates to shed light on the present disagreement of the observed $\sin 2\beta$ between $B \to J/\psi K$ and $B \to \phi K$ decays if it exists.

Furthermore, since the complex phase of right-handed couplings is a new source of CP violation, the rate asym-



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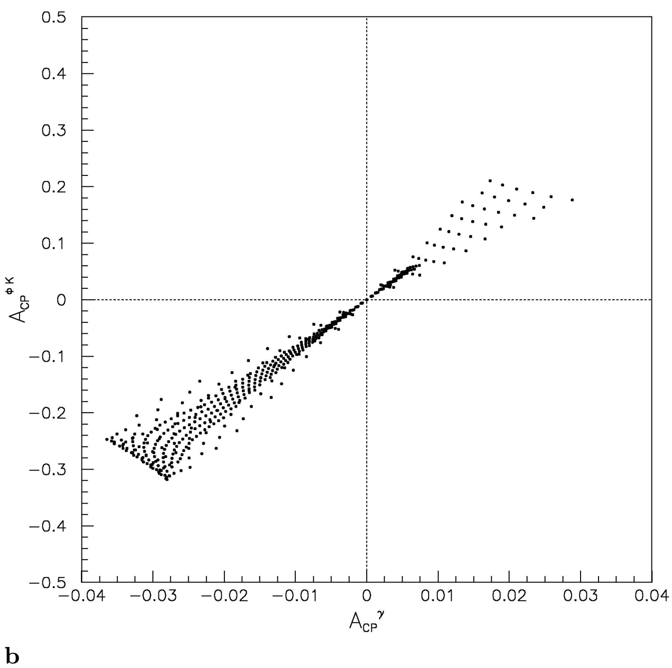

b

Fig. 6. a Correlation of the CP asymmetry of $B \to X_s\gamma$ decay and $\sin 2\beta_{\rm eff}^{\phi K}$ with varying ξ_b as shown in Fig. 1. **b** Correlation of the CP asymmetry of $B \to X_s\gamma$ decay and $A_{CP}^{\phi K}$ with varying ξ_b as shown in Fig. 1

metry indicating a direct CP violation also exists in $B \to \phi K$ decays. This CP asymmetry may be large, up to -30% and may be another strong indication of the right-handed top couplings.

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