CP Violation in $B \to \phi K$ Decay

with Anomalous Right-handed Top Quark Couplings

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Abstract

We explore the 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 the new CP violation and lead to a deviation of the observed weak phase in $B \to \phi K$ decays, which takes account for the present disagreement of the observed $\sin 2\beta$ between $B \to J/\psi K$ and $B \to \phi K$ decays. The direct CP violation is also predicted.

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I. INTRODUCTION

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

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

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 $O(\lambda^2) \sim 5\% [3,4]$. Therefore a sizable disagreement of $\sin 2\beta$ between $B \to \phi K$ 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]

$$\sin 2\beta = 0.734 \pm 0.054,$$

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

The left-right (LR) model based on $\text{SU}(2)_L \times \text{SU}(2)_R \times \text{U}(1)$ gauge group is one of the natural extensions of the SM [7]. In the LR model, the right-handed quark mixing is also an observable as well as the left-handed quark mixing. Without manifest symmetry between left- and right-handed sectors, the right-handed quark mixing is not necessarily 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 left ones, which are suppressed by the heavy mass of extra $W$ boson. The strength of right-handed couplings should be determined by measurements in various phenomena. On the other hand, when the electroweak symmetry is dynamically broken, some nonuniversal 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 Large Hadron Collider (LHC) will allow us to study the structure of 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 direct way. The subdominant channel of the top quark is the CKM nondiagonal decay $t \to sW$ in the SM of which branching ratio is estimated as $\text{Br}(t \to sW) \sim 1.6 \times 10^{-3}$, when $|V_{ts}| = 0.04$ is assumed. Despite the small branching ratio of this channel, the large number of expected top quark production at LHC will enable us to measure the $t \to sW$ process and provide us 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 anomalous right-handed $tbW$ and $tsW$ couplings and their impact on the CP violation in $B \rightarrow \phi K_S$ decay.

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

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

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

This paper is organized as follows: In section II, the effective Hamiltonian formalism with right-handed $tbW$ and $tsW$ couplings is given. In section III, we discuss the constraints on the parameters $\xi_b$ and $\xi_s$ using the radiative $B \rightarrow X_s \gamma$ decay. The analysis on hadronic decays $B \rightarrow J/\psi K$ and $B \rightarrow \phi K$ is presented in section IV to extract corresponding $\sin 2\beta$. Finally we conclude in section V.

II. 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

$$\mathcal{H}_{eff} = \frac{4G_F}{\sqrt{2}} V_{tb} V_{tb}^\ast \left[ \sum_{i=1}^{2} (C_i(\mu)O_i(\mu) + C'_i(\mu)O'_i(\mu)) - \sum_{i=3}^{10} (C_i(\mu)O_i(\mu) + C'_i(\mu)O'_i(\mu)) \right] + \text{H.c.}, \quad (4)$$

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

$$O_1 = (\bar{s}_\alpha c_\beta)_{L} (\bar{c}_\beta b_\alpha)_{L},$$
$$O_2 = (\bar{s}c)_{L}(\bar{c}b)_{L},$$
$$O_3 = (\bar{s}b)_{L} \sum_{q' = u,d,s,c,b} (\bar{q}' q')_{L},$$
$$O_4 = (\bar{s}_\alpha b_\beta)_{L} \sum_{q' = u,d,s,c,b} (\bar{q}' q'_\alpha)_{L},$$
$$O_5 = (\bar{s}b)_{L} \sum_{q' = u,d,s,c,b} (\bar{q}' q')_{R},$$
$$O_6 = (\bar{s}_\alpha b_\beta)_{L} \sum_{q' = u,d,s,c,b} (\bar{q}' q'_\alpha)_{R},$$
$$O_7 = \frac{3}{2}(\bar{s}b)_{L} \sum_{q' = u,d,s,c,b} e_{q'}(\bar{q}' q')_{L},$$
$$O_8 = \frac{3}{2}(\bar{s}_\alpha b_\beta)_{L} \sum_{q' = u,d,s,c,b} e_{q'}(\bar{q}' q'_\alpha)_{L}. $$
\( O_9 = \frac{3}{2} (\bar{s}b)_L \sum_{q'=u,d,s,c,b} e_{q'} (\bar{q}'q')_R, \)
\( O_{10} = \frac{3}{2} (\bar{s}_t b_\beta)_L \sum_{q'=u,d,s,c,b} e_{q'} (\bar{q}'_\beta q'_\alpha)_R, \)
\( O_{11} = \frac{g_s}{16\pi^2} m_b \bar{s}_t \sigma_{\mu\nu} T_{\alpha\beta} b_\beta G^a \mu\nu, \)
\( O_{12} = \frac{e}{16\pi^2} m_b \bar{s}_t \sigma_{\mu\nu} b F^{\mu\nu}, \)

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

Matching the effective Hamiltonian and our model lagrangian of Eq. (2) at \( \mu = m_W \) scale, we have the Wilson coefficients \( C_i (\mu = m_W) \) and \( C'_i (\mu = m_W) \) in the SM:

\[
\begin{align*}
C_1 (m_W) &= \frac{11 \alpha_s (m_W)}{2} 
C_2 (m_W) &= 1 - \frac{11 \alpha_s (m_W)}{6} - \frac{35 \alpha}{18 \pi} 
C_3 (m_W) &= -\frac{\alpha_s (m_W)}{24\pi} E_0 (x_t) + \frac{\alpha}{6\pi \sin^2 \theta_W} [2 B_0 (x_t) + C_0 (x_t)], \\
C_4 (m_W) &= \frac{\alpha_s (m_W)}{8\pi} E_0 (x_t),
C_5 (m_W) &= -\frac{\alpha_s (m_W)}{24\pi} E_0 (x_t),
C_6 (m_W) &= \frac{\alpha_s (m_W)}{8\pi} E_0 (x_t),
C_7 (m_W) &= \frac{\alpha}{6\pi} [4 C_0 (x_t) + D_0 (x_t)],
C_9 (m_W) &= \frac{\alpha}{6\pi} [4 C_0 (x_t) + D_0 (x_t) + \frac{1}{\sin^2 \theta_W} (10 B_0 (x_t) - 4 C_0 (x_t))],
C_8 (m_W) &= C_{10} (m_W) = 0,
C_{11} (m_W) &= G (x_t),
C_{12} (m_W) &= F (x_t),
C'_i (m_W) &= 0, \quad i = 1, \cdots, 12,
\end{align*}
\]

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 explicit forms are given in Refs. [13,14]. Turning on the right-handed \( \bar{t}bW \) and \( \bar{t}sW \) couplings, we have the modification of loop functions in the Wilson coefficients \( C_i \),

\[
\begin{align*}
D_0 (x_t) &\to D_0 (x_t) + \xi_b \frac{m_b}{m_t} D_R (x_t), \\
E_0 (x_t) &\to E_0 (x_t) + \xi_b \frac{m_b}{m_t} E_R (x_t), \\
F (x_t) &\to F (x_t) + \xi_t \frac{m_t}{m_b} F_R (x_t), \\
G (x_t) &\to G (x_t) + \xi_t \frac{m_t}{m_b} G_R (x_t),
\end{align*}
\]
and also have the new Wilson coefficients $C'_i$:

$$
C'_3(m_W) = -\frac{\alpha_s(m_W)}{24\pi} \xi_s \frac{m_b}{m_t} E_R(x_t),
$$
$$
C'_4(m_W) = \frac{\alpha_s(m_W)}{8\pi} \xi_s \frac{m_b}{m_t} E_R(x_t),
$$
$$
C'_5(m_W) = -\frac{\alpha_s(m_W)}{24\pi} \xi_s \frac{m_b}{m_t} E_R(x_t),
$$
$$
C'_6(m_W) = \frac{\alpha_s(m_W)}{8\pi} \xi_s \frac{m_b}{m_t} E_R(x_t),
$$
$$
C'_7(m_W) = \frac{\alpha}{6\pi} \xi_s \frac{m_b}{m_t} D_R(x_t),
$$
$$
C'_8(m_W) = \frac{\alpha}{6\pi} \xi_s \frac{m_b}{m_t} D_R(x_t),
$$
$$
C'_{11}(m_W) = \xi_s \frac{m_l}{m_b} G_R(x_t),
$$
$$
C'_{12}(m_W) = \xi_s \frac{m_i}{m_b} F_R(x_t),
$$
$$
C'_i(m_W) = 0, \quad (i = 1, 2, 8, 10),
$$

(8)

where the new loop functions are given by

$$
D_R(x) = \frac{x(59 - 38x + 25x^2 + 2x^3)}{36(x - 1)^4} + \frac{2(x + 1)}{3(x - 1)^5} \ln x + \frac{x^2}{2(x - 1)^4} \ln x,
$$
$$
E_R(x) = \frac{x(-116 + 165x - 114x^2 + 29x^3)}{18(x - 1)^4} + \frac{2 + 3x + x^2}{3(x - 1)^5} \ln x,
$$
$$
F_R(x) = \frac{-20 + 31x - 5x^2}{12(x - 1)^2} + \frac{x(2 - 3x)}{2(x - 1)^3} \ln x,
$$
$$
G_R(x) = -\frac{4 + x + x^2}{4(x - 1)^2} + \frac{3x}{2(x - 1)^3} \ln x,
$$

(9)

where our new loop functions $F_R(x), G_R(x)$ and $D_R(x)$ agree with those in Ref. [15,16] and $E_R(x)$ is the first calculation. Note that the $O(\xi)$ terms of $Z$-penguin diagram are suppressed by the heavy mass of $Z$-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 two currents are different and one current is proportional to the fermion momentum and the other current proportional to the fermion mass. It indicates that the integrand is always an odd function and the loop integral vanishes. Therefore the $O(\xi)$ terms of the box diagram do not exist and the leading contribution is of $\xi^2$ order. Hence, we also ignore the box contribution. As a consequence, the contribution of order $O(\xi)$ comes only through the $\gamma$-penguin and gluon penguin diagrams. Actually the contributions of $O(\xi)$ operator to hadronic decays are very small and we neglect it in the numerical analysis.

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

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

5
is governed by a $24 \times 24$ anomalous dimension matrix $\gamma$. Since the strong interaction preserves chirality, the operators $O_i$ and $O'_i$ are evolved separately without mixing between them. Thus the $24 \times 24$ anomalous dimension matrix $\gamma$ is decomposed into two identical $12 \times 12$ matrices $\gamma_0$ given in the SM. The $12 \times 12$ anomalous dimension matrix $\gamma_0$ can be found in Ref. [13,17,18]. The evolved Wilson coefficients $C_i(\mu)$ are expressed in terms of the initial conditions of Eq. (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 Ref. [17,18]. The matrix elements of 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 Ref. [19–21]. Then the Hamiltonian is expressed in terms of effective Wilson coefficients and tree level matrix elements.

III. $B \to X_S \gamma$ CONSTRAINTS

Before the analysis on $\sin 2\beta$, we consider the radiative $B \to X_S \gamma$ decay to constrain the model. This channel has already been observed experimentally and more precise measurement will be obtained from the accumulation of 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 the stringent limits on $\xi_b$ and $\xi_s$ are yielded from the measurement of the $B \to X_S \gamma$ decay [15,9]. We present the updated constraints on 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

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

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 as

$$A_{CP}(B \to X_S \gamma) = \frac{\Gamma(B \to X_S \gamma) - \Gamma(B \to \bar{X}_S \gamma)}{\Gamma(B \to X_S \gamma) + \Gamma(B \to \bar{X}_S \gamma)}$$

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]

$$A_{CP}(B \to X_S \gamma) = (-0.079 \pm 0.108 \pm 0.022)(1.0 \pm 0.030),$$

where the first error is statistical, the second is additive systematic over the various $b \to s \gamma$ decay modes, and the third is multiplicative systematic. Note that the present measurement of $A_{CP}$ is still consistent with 0. The complex anomalous $tbW$ coupling contributes to the CP asymmetry through the interference terms of $O_{11}$ and $O_{12}$ operators such as $\delta A_{CP} \sim a_1 ImC_2C_{12}^* + a_2 ImC_{11}C_{12}^* + a_3 ImC_2C_{11}^*$, which provides an additional test on $\xi_b$, independent of the branching ratio. On the contrary, the $tsW$ coupling does not contribute to $A_{CP}$ at this level.
The explicit expressions of branching ratio and the CP asymmetry are presented in Refs. [26] and [27] in terms of the evolved Wilson coefficients at $\mu = m_b$ scale. With the measured values of Eqs. (9) and (10), we obtain the constraints on $\xi_b$ at 2σ C.L. as

$$-0.002 < Re\xi_b + 22|\xi_b|^2 < 0.0033,$$

$$-0.299 < \frac{0.27 Im\xi_b}{0.095 + 12.54 Re\xi_b + 414.23|\xi_b|^2} < 0.141,$$

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

$$|\xi_s| < 0.012,$$

from the branching ratio alone [15].

IV. HADRONIC DECAYS

A. $B \to J/\psi K$

The $B \to J/\psi K$ decays are dominated by the tree-level $b \to c\bar{s}s$ 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 $\beta$. The CP asymmetries in $B \to J/\psi K_S$ decays given in Eq. (2), $\sin 2\beta_{\text{eff}} = 0.734 \pm 0.054$, agrees 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 suppression factor is estimated of order $< 10^{-4}$. Thus new physics effect on 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 $\mathcal{O}(\xi_b^2)$,

$$M_{12} = M_{12}^{SM} \left( 1 + \xi_b^2 \frac{S_R(x_t)}{S_0(x_t)} \frac{(b\bar{P}_Ld)(\bar{b}\bar{P}_Ld)}{(b\gamma_{\mu}\bar{P}_Ld)(b\gamma_{\mu}\bar{P}_Ld)} \right),$$

where the new loop function $S_R(x)$ is given by

$$S_R(x) = \frac{x(x^2 - 2x + 6)}{(1 - x)^2} + \frac{x(x + 2)(x^2 - x + 2)}{(1 - x)^3} \ln x,$$

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

$$\sin 2\beta_{\text{eff}} = \sin 2\beta + 4.3|\xi_b|^2 \sin 2\varphi,$$
where \( \xi_b = |\xi_b| e^{i\varphi} \) and

\[
\frac{\langle B^0| (\bar{b}pLd)(\bar{b}pLd)|B^0 \rangle}{\langle B^0| (b\gamma_\mu P_Ld)(b\gamma^-\mu P_Ld)|B^0 \rangle} \approx \frac{3}{4} \left( \frac{m_B}{m_b} \right)^2.
\]

With the allowed parameter set of Fig. 1, the second term of Eq. (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\( \beta_{\text{eff}} \) 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.

**B. \( B \to \phi K_S \)**

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

\[
A_{CP}^{\phi K} = -0.56 \pm 0.43, \quad \sin 2\beta_{\text{eff}}^{\phi K} = -0.39 \pm 0.41,
\] (17)

where \( A_{CP}^{\phi K} \) is the CP violating asymmetry defined as \( A_{CP}^{\phi K} \equiv \frac{\Gamma(B \to \phi K) - \Gamma(\bar{B} \to \phi \bar{K})}{\Gamma(B \to \phi K) + \Gamma(\bar{B} \to \phi \bar{K})} \) and \( \sin 2\beta_{\text{eff}}^{\phi K} \) 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 by the penguin loop makes the new effect of \( O_{11} \) 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 \( B \to \phi K_S \) 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 \( \bar{t}bW \) and \( \bar{t}sW \) couplings and compare \( A_{CP}^{\phi K} \) and \( \sin 2\beta_{\text{eff}}^{\phi K} \) with the experiment. On the other hand, the electroweak penguin operators also give 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

\[
\langle P(p')|V_{\mu}|B(p)\rangle = \left[ (p + p')_{\mu} - \frac{m_B^2 - m_p^2}{q^2} q_{\mu} \right] F_1^p(q^2) + \frac{m_B^2 - m_p^2}{q^2} q_{\mu} F_0^p(q^2),
\]

\[
\langle 0|A_{\mu}|P(p)\rangle = i f_{Pp} p_{\mu},
\]

\[
\langle 0|V_{\mu}|V(p)\rangle = f_V m_V \epsilon_{\mu},
\] (18)

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

\[
\mathcal{A}(B^0 \to \phi K^0) = \frac{G_F}{\sqrt{2}} V_{tb}^* V_{ts} \left[ a_3 + a_4 + a_5 - \frac{1}{2} (a_7 + a_9 + a_{10}) \right] 2 f_{\phi} m_\phi (e^* \cdot p_B) F_1^K + A_{11}^{\phi K},
\] (19)
where \(a_{2i-1} = C_{2i-1} + C_{2i}/N_c\), \(a_{2i} = C_{2i} + C_{2i-1}/N_c\). Contribution of the chromo-magnetic operator \(A_{11}^{\phi K}\) is given by [28,29]

\[
A_{11}^{\phi K} = \langle \phi K^0 | (H_{11} + H'_{11}) | B^0 \rangle, \\
= \frac{G_F \alpha_s(q^2)}{\sqrt{2}} \frac{N_c^2 - 1}{N_c^2} \frac{N_c}{m_b} f_\phi m_\phi (e^* p_B) (C_{11} + C'_{11}) (F_1^{\phi K} X + F_0^{\phi K} Y),
\]

with

\[
X = 4m_b + 5m_s + 3m_s \left( \frac{m_B^2 - m_K^2}{m_\phi^2} \right) - \left( \frac{3m_B^2 - 3m_K^2 + m_\phi^2}{8m_b} \right) \left( 1 + \frac{m_B^2 - m_K^2}{m_\phi^2} \right),
\]

\[
Y = \frac{3}{2} \left( \frac{m_B^2 - m_K^2}{m_b - m_s} \right) + \left( \frac{m_B^2 - m_K^2}{m_\phi^2} \right) \left( \frac{3m_B^2 - 3m_K^2 + m_\phi^2}{8m_b} \right) - 3m_s,
\]

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. The early calculation was performed in the framework of a quark model [30]. We can set \(F_0 = F_1\) close to the point \(q^2 = 0\) [31] and assume the 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 four-quark operators are doubly suppressed by both \(m_b/m_t\) and \(\xi_{b,s}\), while the effects on 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 notify that \(C_i\) in the Eqs. (19) and (20) are the effective Wilson coefficients absorbing the 1-loop correction to the hadronic matrix elements and they involve the strong phases.

Since the four-quark operator contribution in the first term in Eq. (19) involves a strong phase, the new phase of \(\phi K\) 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 \(\lambda\) defined by

\[
\lambda = \frac{M_{12} \bar{A}}{M_{12} A},
\]

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

\[
a_{\phi K}(t) = \frac{\Gamma(B_{\text{phys}}^0(t) \to \phi K^0) - \Gamma(B_{\text{phys}}^0(t) \to \phi \bar{K}^0)}{\Gamma(B_{\text{phys}}^0(t) \to \phi K^0) + \Gamma(B_{\text{phys}}^0(t) \to \phi \bar{K}^0)},
\]

\[
= C_{\phi K} \cos \Delta m_B t + S_{\phi K} \sin \Delta m_B t,
\]

where the coefficients are

\[
C_{\phi K} = \frac{1 - |\lambda|^2}{1 + |\lambda|^2} \equiv -A_{CP}^{\phi K},
\]

\[
S_{\phi K} = \frac{2Im\lambda}{1 + |\lambda|^2} \equiv \sin 2\beta_{\text{eff}}^{\phi K}.
\]

Note that the hadronic uncertainty is cancelled in \(\lambda\) and the CP violating observables \(A_{CP}^{\phi K}\) and \(\sin 2\beta_{\text{eff}}^{\phi K}\) are free from the hadronic uncertainty. We can express the parameter \(\lambda\) by
\[ \lambda = \lambda^{\text{SM}} \left( 1 + 21.84 \ e^{-i \delta} \ \xi_q \right) \approx \lambda^{\text{SM}} (1 + i \ 43.7 \ |\xi_q| \ e^{-i \delta} \ \sin \varphi_q) \]  

where \( \varphi_q \) is the phase of \( \xi_q \), and \( \delta \) the strong phase introduced by the one loop correction to the matrix elements. For \( \lambda^{\text{SM}} \equiv e^{i \beta^{\text{SM}}} \), we will use the measured value given in Eq. (2). With this expression, we can write the CP asymmetries as

\[ A_{CP}^{\phi K} = 23.3 \ |\xi_q| \ \sin \varphi_q, \]

\[ \sin 2\beta_{\text{eff}}^{\phi K} = \sin 2\beta + 52.2 \ |\xi_q| \ \sin \varphi_q, \]  

(26)

where \( \delta = 2.58 \) in our calculation. Note that the second expression of Eq. (25) is no more valid for the maximal value of \( |\xi_b| \sim 0.04 \) and so are the above expressions of \( A_{CP}^{\phi K} \) and \( \sin 2\beta_{\text{eff}}^{\phi K} \).

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

\[ -0.34 < A_{CP}^{\phi K} < 0.22, \]

\[ -0.10 < \sin 2\beta_{\text{eff}}^{\phi K} < 0.96. \]  

(27)

As shown in the previous section, the effect of the right-handed top couplings on \( B - \bar{B} \) mixing sector is safely neglected for the evaluation of \( \sin 2\beta_{\text{eff}}^{\phi K} \). In Fig. 2, the correlation of \( \sin 2\beta_{\text{eff}}^{\phi K} \) 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 experiment ascertain that the \( \sin 2\beta_{\text{eff}}^{\phi K} \) 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

\[ -0.28 < A_{CP}^{\phi K} < 0.28, \]

\[ 0.23 < \sin 2\beta_{\text{eff}}^{\phi K} < 0.94, \]  

(28)

and their correlation is shown in Fig. 3. We also find that the large CP asymmetry \((-10 \%)\) is possible with \( \xi_s \) even if \( \sin 2\beta_{\text{eff}}^{\phi K} \) agrees with the SM prediction.

We also calculate the branching ratio:

\[ Br(B \to \phi K) = \frac{\tau_B}{16\pi} \frac{\lambda(m_B^2, m_{\phi}^2, m_K^2)}{m_B^3} |A|^2, \]  

where \( \lambda(x, y, z) = (x^2 + y^2 + z^2 - 2xy - 2yz - 2zx)^{1/2} \) and \( \tau_B \) is the lifetime of \( B \) meson. Figure 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 \( B^0 \to \phi K^0 \) decay read:

\[ Br(B^0 \to \phi K^0) = (5.4^{+3.7}_{-2.7} \pm 0.7) \times 10^{-6} < 12.3 \times 10^{-6} \quad \text{CLEO}, \]

\[ = (8.1^{+3.1}_{-2.5} \pm 0.8) \times 10^{-6} \quad \text{BaBar}, \]

\[ = (8.7^{+3.8}_{-3.0} \pm 1.5) \times 10^{-6} \quad \text{Belle}, \]  

(30)

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 the Fig. 4 and 5. In our evaluation, the SM value is $(1.9 - 4.0) \times 10^{-6}$ close to the prediction of Ref. [38]. We can find that the negative $\sin 2\beta_{\text{eff}}$ consistent with the BaBar cross section is possible with anomalous $\bar{t}bW$ coupling but we do not expect such a solution with anomalous $\bar{t}sW$ coupling.

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

V. 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)$_L \times$SU(2)$_R \times$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 $\xi_q$, 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 the negative $\sin 2\beta$ is possible with the anomalous $\bar{t}bW$ coupling, as is consistent with the recent measurements. Note that the $\sin 2\beta$ with the anomalous $\bar{t}sW$ coupling is also shifted but still positive due to more strict bound on $|\xi_s|$ than that on $|\xi_b|$. In conclusion, the right-handed top couplings are one of the good candidates of 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 asymmetry indicating a direct CP violation also exists in $B \to \phi K$ decays. This CP asymmetry may be large, up to $-30\%$ and another strong indication of the right-handed top couplings.

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FIGURES
FIG. 1. Allowed parameter set \((Re\xi_b, Im\xi_b)\) under the constraints by the branching ratio and CP asymmetry in \(B \to X_s \gamma\) decay.
FIG. 2. Correlation of $\sin 2\beta_{\text{eff}}$ and $A_{CP}^{\phi K}$ with varying $\xi_b$ shown in Fig. 1.
FIG. 3. Correlation of $\sin 2\beta_{\text{eff}}$ and $A_{\text{CP}}^{\phi K}$ with varying $\xi_s$ under the constraint of Eq. (11).
FIG. 4. (a) Correlation of the branching ratio of $B \to \phi K$ decay and $\sin 2\beta_{\phi K}$ with varying $\xi_b$ shown in Fig. 1. (b) Correlation of the branching ratio of $B \to \phi K$ decay and $A_{CP}^{\phi K}$ with varying $\xi_b$ shown in Fig. 1.
FIG. 5. (a) Correlation of the branching ratio of $B \to \phi K$ decay and $\sin 2\beta_{\text{eff}}^{\phi K}$ with varying $\xi_s$ under the constraint of Eq. (11). (b) Correlation of the branching ratio of $B \to \phi K$ decay and $A_{CP}^{\phi K}$ with varying $\xi_s$ under the constraint of Eq. (11).
FIG. 6. (a) Correlation of the CP asymmetry of $B \rightarrow X_s \gamma$ decay and $\sin 2\beta_{\text{eff}}$ with varying $\xi_b$ shown in Fig. 1. (b) Correlation of the CP asymmetry of $B \rightarrow X_s \gamma$ decay and $A_{CP}^{\phi K}$ with varying $\xi_b$ shown in Fig. 1.