CP asymmetries in penguin-induced *B* decays in general left-right models

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We study *CP* asymmetries in penguin-induced $b\to s\bar{s}s$ decays in general left-right models without imposing manifest or pseudomanifest left-right symmetry. Using the effective Hamiltonian approach, we evaluate *CP* asymmetries in $B^{\pm}\to\phi K^{(*)\pm}$ decays as well as mixing induced B meson decays $B\to J/\psi K_S$ and $B\to\phi K_S$ decays. Based on recent measurements revealing large *CP* violation, we show that a nonmanifest type model is more favored than a manifest or pseudomanifest type.

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I. INTRODUCTION

One of the major goals of present experiments in B physics is the study of *CP* violation which may reside in the quark flavor mixing described by the Cabibbo-Kobayashi-Maskawa (CKM) matrix in the standard $SU(2)_L\times U(1)$ model (SM). Since there is one complex phase in the CKM matrix, the sizes and patterns of *CP* violation in various decay modes in the SM are in principle expressed through this single parameter [1]. But the present experimental results with large *CP* violation effects in the B meson system are not simply explained with this single parameter under the minimal SM framework [2]. For instance, the *CP* asymmetries in mixing induced B meson decays are characterized by a *CP* angle β which is a phase of the CKM matrix element V_{td}, and the observed world average value of $\sin 2\beta$ in $B\to J/\psi K_S$ ($b\to c\bar{c}s$) decays is given by

$$\sin 2\beta_{J/\psi K_S}=0.734\pm 0.054. \tag{1}$$

In addition, this *CP* angle β is recently measured by BABAR and Belle in $B\to\phi K_S$ ($b\to s\bar{s}s$) decays [3], and their average value is

$$\sin 2\beta_{\phi K_S}=-0.39\pm 0.41. \tag{2}$$

In the SM, however, the *CP* asymmetry in $B\to\phi K_S$ decays is expected to be very close to that in $B\to J/\psi K_S$ decays [4]. Admitting that the statistical error of those experimental data is still too large to confirm the data and justify any theory, a 2.7σ deviation between $\sin 2\beta_{J/\psi K_S}$ and $\sin 2\beta_{\phi K_S}$ may give a clue of new physics (NP) effects in B decays. If so, other inclusive $b\to s\bar{s}s$ dominated B decays such as $B^{\pm}\to\phi K^{(*)\pm}$ decays might receive the same contribution from the NP.

In a recent paper [5], we have investigated the mixing induced *CP* asymmetry in $B\to J/\psi K_S$ decays in the general left-right model (LRM) with group $SU(2)_L\times SU(2)_R\times U(1)$ since it is one of the simplest extensions of the SM gauge group as a complement of the purely left-handed nature of the SM [6]. Because of the extended group $SU(2)_R$ in the LRM there are new neutral and charged gauge bosons, Z_R and W_R as well as a right-handed gauge coupling, g_R. After spontaneous symmetry breaking, the gauge eigenstates W_R mix with W_L to form the mass eigenstates W and W' with masses M_W and $M_{W'}$, respectively. The W_L-W_R mixing angle ξ and the ratio ζ of M_W^2 to $M_{W'}^2$ are restricted by a number of low-energy phenomenological constraints along with the right-handed mass mixing matrix elements. From the limits on deviations of muon decay parameters from the V-A prediction, the lower bound on $M_{W'}$ can be obtained as follows [7]:

$$\zeta_g<0.033 \quad \text{or} \quad M_{W'}>(g_R/g_L)\times 440\ \text{GeV}, \tag{3}$$

where $\zeta_g\equiv g_R^2M_W^2/g_L^2M_{W'}^2$. Previously, stronger limits of the mass $M_{W'}$ as well as the mixing angle ξ were presented by many authors experimentally [8] and theoretically [9] assuming manifest ($V^R=V^L$) or pseudomanifest ($V^R=V^{L*}K$) left-right symmetry ($g_L=g_R$), where V^L and V^R are the left-and right-handed quark mixing matrices, respectively, and K is a diagonal phase matrix [10]. But, in general, the form of V^R is not necessarily restricted to manifest or pseudomanifest symmetric types, so the W_R mass limit can be lowered to approximately 300 GeV by taking the following forms of V^R [11]:

$$V_I^R=\begin{pmatrix} e^{i\omega} & \sim 0 & \sim 0 \\ \sim 0 & c_Re^{i\alpha_1} & s_Re^{i\alpha_2} \\ \sim 0 & -s_Re^{i\alpha_3} & c_Re^{i\alpha_4}\end{pmatrix},$$

$$V_{II}^R=\begin{pmatrix} \sim 0 & e^{i\omega} & \sim 0 \\ c_Re^{i\alpha_1} & \sim 0 & s_Re^{i\alpha_2} \\ -s_Re^{i\alpha_3} & \sim 0 & c_Re^{i\alpha_4}\end{pmatrix}, \tag{4}$$

where $c_R(s_R)\equiv\cos\theta_R(\sin\theta_R)$ $(0°\leqslant\theta_R\leqslant 90°)$. Here the matrix elements indicated as ~ 0 may be $\lesssim 10^{-2}$ and unitarity requires $\alpha_1+\alpha_4=\alpha_2+\alpha_3$. From the $b\to c$ semileptonic decays of the B mesons, we can get an approximate bound $\xi_g\sin\theta_R\lesssim 0.013$ by assuming $|V_{cb}^L|\approx 0.04$ [12], where $\xi_g\equiv(g_R/g_L)\xi$.[1] This new parameter ξ_g is in general smaller than the charged gauge boson mass ratio ζ_g in the general

[1]In Ref. [5], ξ_g is defined as $(g_L/g_R)\xi$ unlike this paper so that the mistakenly written bound $\xi_g\sin\theta_R\lesssim 0.013$ should read $(g_R/g_L)\xi\sin\theta_R\lesssim 0.013$.

LRM [5,9]. In a similar way to the charged gauge bosons, the neutral gauge bosons mix each other [13]. But we do not present them here because Z_R contribution to penguin-induced B decays is negligible. Also, due to gauge invariance, tree-level flavor-changing neutral Higgs bosons with masses M_H enter into our theory [14]. However, we also neglect their contributions by assuming $M_H \gg M_{W'}$.

The CP asymmetry in the penguin-induced $B\to\phi K_S$ decays was also studied earlier in the pseudomanifest left-right symmetry model in Ref. [15]. In this case, the right-handed current contribution to $B\bar{B}$ mixing is suppressed by ζ so that the NP effect only arises in the magnetic penguin since the suppression by ξ is offset by a large factor m_t/m_b arising in the virtual top quark loop [16]. However, in the nonmanifest LRM, ζ terms in $B\bar{B}$ mixing and absorptive part of the decay amplitudes become important due to the possible enhancement of V^R elements so that the right-handed current contribution to the corresponding CP asymmetry is more enhanced. In this paper, as a continuation of our previous work, we will explicitly evaluate the possible right-handed current contribution to CP asymmetry in $B^\pm\to\phi K^{(*)\pm}$ decays as well as in $B\to\phi K_S$ decays in the general LRM related to recent measurements, and show that CP asymmetries in those decays can be large enough to probe the existence of the right-handed current using the effective Hamiltonian approach. After reviewing the structure of the effective Hamiltonian in the general LRM in Sec. II, we will discuss CP asymmetries in the several $b\to s\bar{s}s$ dominated B decays in Sec. III in detail.

II. EFFECTIVE HAMILTONIAN

The low-energy effects of the full theory can be described by the effective Hamiltonian approach in order to include QCD effects systematically. The low-energy effective Hamiltonian calculated within the framework of the operator product expansion (OPE) has a finite number of operators in a given order, which is dependent upon the structure of the model. In the LRM, the low energy effective Hamiltonian at the energy scale μ for $\Delta B=1$ and $\Delta S=1$ transition has the following form:

$$\mathcal{H}_{eff}=\frac{G_F}{\sqrt{2}}\left[\sum_{\substack{i=1,2\\ q=u,c}}\lambda_q^{LL}C_i^qO_i^q-\lambda_t^{LL}\left(\sum_{i=3}^{12}C_iO_i+C_7^\gamma O_7^\gamma+C_8^GO_8^G\right)\right]+(C_iO_i\to C_i'O_i'), \tag{5}$$

where $\lambda_q^{AB}\equiv V_{qs}^{A*}V_{qb}^B$, $O_{1,2}$ are the standard current-current operators, O_3-O_{10} are the standard penguin operators, and O_7^γ and O_8^G are the standard photonic and gluonic magnetic operators, respectively, which can be found in Ref. [17]. Since we have additional $SU(2)_R$ group in the LRM, the operator basis is doubled by O_i' which are the chiral conjugates of O_i. Also new operators $O_{11,12}$ and $O_{11,12}'$ arise with mixed chiral structure of $O_{1,2}$ and $O_{1,2}'$ [16].

In order to calculate the Wilson coefficients $C_i(\mu)$, we first calculate them at $\mu=M_W$ scale. After performing a straightforward matching computation, we find the Wilson coefficients at W scale neglecting the u-quark mass:

$$\begin{aligned}
&C_2^q(M_W)=1,\quad C_2^{q\prime}(M_W)=\zeta_g\lambda_q^{RR}/\lambda_q^{LL},\\
&C_7^\gamma(M_W)=F(x_t^2)+A^{tb}\tilde{F}(x_t^2),\\
&C_7^{\gamma\prime}(M_W)=A^{ts*}\tilde{F}(x_t^2),\\
&C_8^G(M_W)=G(x_t^2)+A^{tb}\tilde{G}(x_t^2),\\
&C_8^{G\prime}(M_W)=A^{ts*}\tilde{G}(x_t^2),
\end{aligned} \tag{6}$$

where

$$x_q=\frac{m_q}{M_W}\ (q=u,c,t),\quad A^{tD}=\xi_g\frac{m_t}{m_b}\frac{V_{tD}^R}{V_{tD}^L}e^{i\alpha_\circ}\ (D=b,s), \tag{7}$$

and $\alpha_\circ$ is a CP phase residing in the vacuum expectation values, which can be absorbed in α_i in Eq. (4) by redefining $\alpha_i+\alpha_\circ\to\alpha_i$. All other coefficients vanish. In Eq. (6), the explicit forms of the functions $F(x_t)$, $\tilde{F}(x_t)$, $G(x_t)$, and $\tilde{G}(x_t)$ are given in Ref. [16], and the terms proportional to ξ_g and ζ_g in the magnetic coefficients are neglected except the contribution coming from the virtual t quark which gives m_t/m_b enhancement. Also the term proportional to ζ_g in the tree-level coefficient C_2' is not neglected because $\zeta_g \gg \xi_g$ and there is possible enhancement by the ratio of CKM angles ($\lambda_q^{RR}/\lambda_q^{LL}$) in the nonmanifest LRM.

The coefficients $C_i(\mu)$ at the scale $\mu=m_b$ can be obtained by evolving the coefficients $C_i(M_W)$ with the 28 $\times$28 anomalous dimension matrix applying the usual renormalization group procedure. Since the strong interaction preserves chirality, the 28$\times$28 anomalous dimensional matrix decomposes into two identical 14$\times$14 blocks. The SM 12 $\times$12 submatrix describing the mixing among O_1-O_{10}, O_7^γ, and O_8^G can be found in Ref. [18], and the explicit form of the remaining 4$\times$4 matrix describing the mixing among $O_{11,12}$, O_7^γ, and O_8^G, which partially overlaps with the SM 12$\times$12 submatrix, can be found in Ref. [16]. The low energy Wilson coefficients at the scale $\mu=m_b$ in the LL approximation are then given by

$$C_i(m_b)=\sum_{j,k}(S^{-1})_{ij}(\eta^{3\lambda_j/23})S_{jk}C_k(M_W), \tag{8}$$

where the λ_j's in the exponent of $\eta=\alpha_s(M_W)/\alpha_s(m_b)$ are the eigenvalues of the anomalous dimension matrix over $g^2/16\pi^2$ and the matrix S contains the corresponding eigenvectors. The result for the photonic and gluonic magnetic coefficients are calculated in Ref. [16] and in Ref. [15], respectively, and the rest of them related to our analysis can be

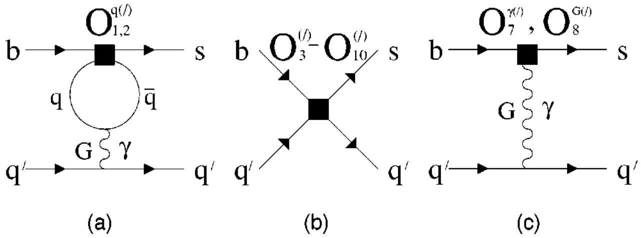


FIG. 1. Diagrams for penguin-induced $b\to s\overline{q'q'}$ decays.

found in Ref. [17].[2] Therefore we do not repeat them here, and lead the reader to the original papers. For 5 flavors, we have the following numerical values of $C_i(m_b)$ in LL precision using $\Lambda_{\overline{MS}}=225$ MeV, $m_b=4.4$ GeV, and $m_t=170$ GeV:[3]

$$C_1^q=-0.308,\quad C_1^{q\prime}=C_1^q\zeta_g\lambda_q^{RR}/\lambda_q^{LL},$$

$$C_2^q=1.144,\quad C_2^{q\prime}=C_2^q\zeta_g\lambda_q^{RR}/\lambda_q^{LL},$$

$$C_3=0.014,\quad C_4=-0.030,\quad C_5=0.009,\quad C_6=-0.038,$$

$$C_7=0.045\alpha,\quad C_8=0.048\alpha,$$

$$C_9=-1.280\alpha,\quad C_{10}=0.328\alpha, \tag{9}$$

$$C_7^\gamma=-0.317-0.546A^{tb},\quad C_7^{\gamma\prime}=-0.546A^{ts*},$$

$$C_8^G=-0.150-0.241A^{tb},\quad C_8^{G\prime}=-0.241A^{ts*}.$$

Note that $C_3'-C_{10}'$ are negligible comparing to $C_7^{\gamma\prime}$ and $C_8^{G\prime}$ whereas $C_{1,2}'$ are not. We will show that $C_{1,2}'$ are important to the absorptive parts in penguin-dominated B decays in the next section.

III. *CP* VIOLATING ASYMMETRIES

A. Charged *B* meson decays

For charged B meson decays, the nonzero *CP* violating asymmetry defined as

$$A_{CP}=\frac{\Gamma(B^+\to f^+)-\Gamma(B^-\to f^-)}{\Gamma(B^+\to f^+)+\Gamma(B^-\to f^-)} \tag{10}$$

originates from the superposition of *CP*-odd(violating) phases introduced by CKM matrix elements and *CP*-even(conserving) phases arising from the absorptive part of the amplitudes. Since we have obtained the relevant effective Hamiltonian in Sec. II, it is quite straightforward to calculate the partial decay rates and *CP* asymmetries in $b\to s\overline{s}s$ decays. These decays are governed by three different types of penguin diagrams shown in Fig. 1. The absorptive part of the amplitudes arises at $O(\alpha_s)$ from the one-loop penguin diagrams with insertions of the operators $O_{1,2}^{(\prime)}$ shown in Fig. 1(a). The detailed calculation of the one-loop penguin matrix element of the operators $O_{1,2}$ in the SM is in Ref. [20] so that we can be very brief. The renormalized matrix elements of the operators $O_{1,2}^{(\prime)}$ in the LL approximation are given by

$$\langle O_1^{q(\prime)}\rangle^{\rm peng}=\frac{\alpha}{3\pi}\mathcal{I}(m_q,k,m_b)\langle P_\gamma^{(\prime)}\rangle,$$

$$\langle O_2^{q(\prime)}\rangle^{\rm peng}=\frac{\alpha_s(m_b)}{8\pi}\mathcal{I}(m_q,k,m_b)\times\left(\langle P_G^{(\prime)}\rangle+\frac{8}{9}\frac{\alpha}{\alpha_s(m_b)}\langle P_\gamma^{(\prime)}\rangle\right), \tag{11}$$

where

$$P_G^{(\prime)}=O_4^{(\prime)}+O_6^{(\prime)}-\frac{1}{N_c}(O_3^{(\prime)}+O_5^{(\prime)}),$$

$$P_\gamma^{(\prime)}=O_7^{(\prime)}+O_9^{(\prime)}\quad(N_c=3), \tag{12}$$

and

$$\mathcal{I}(m,k,\mu)=4\int_0^1 dx\,x(1-x)\ln\left[\frac{m^2-k^2x(1-x)}{\mu^2}\right], \tag{13}$$

and where k is the momentum transferred by the gluon to the $(s,\overline{s})$ pair. As one can see from Eq. (13), different *CP*-even phases arise from the imaginary parts of the functions $\mathcal{I}(m_u,k,\mu)$ and $\mathcal{I}(m_c,k,\mu)$. On the other hand, the penguin operators O_3-O_{10} contribute to only the dispersive parts of the amplitudes and give tree-level penguin transition amplitudes shown in Fig. 1(b). Also, as shown in Fig. 1(c), we should include the tree-level diagram associated with the magnetic operators $O_7^{\gamma(\prime)}$ and $O_8^{G(\prime)}$ to the dispersive part of the amplitude. Using the factorization approximation [21], we use the following parametrization:

$$\langle O_7^{\gamma(\prime)}\rangle^{\rm peng}=-\frac{\alpha}{3\pi}\frac{m_b^2}{k^2}\langle P_\gamma^{(\prime)}\rangle,$$

$$\langle O_8^{G(\prime)}\rangle^{\rm peng}=-\frac{\alpha_s}{4\pi}\frac{m_b^2}{k^2}\langle P_G^{(\prime)}\rangle. \tag{14}$$

Here k^2 is expected to be typically in the range $m_b^2/4\leqslant k^2\leqslant m_b^2/2$ [22]. We will use $k^2=m_b^2/2$ for our numerical analysis.

[2] Although QCD correction factors in $C_{1,2}'$ are different from those in $C_{1,2}$ in general [19], we use an approximation $\alpha_s(M_{W'})\simeq\alpha_s(M_W)$ for simplicity, which will not change our result.

[3] The numbers we obtained for $C_7^{\gamma(\prime)}$ and $C_8^{G(\prime)}$ are slightly different from those in Ref. [15] because they used $m_t/m_b=60$.

Now we are ready to consider $B^{\pm}\to\phi K^{\pm}$ decays explicitly. Since the axial-vector parts of the operators do not contribute to the transition amplitudes in these decays we can simply use $\langle O_i\rangle=\langle O_i'\rangle$ with the help of the vacuum-insertion method [23]. Combining all operators, we obtain the following transition amplitude using the unitarity relation $\Sigma_{q=u,c,t}\lambda_q=0$:

$$\begin{aligned}\mathcal{A}(B^-\to\phi K^-)=\frac{G_F}{\sqrt{2}}\sum_{q=u,c}\lambda_q^{LL}\Bigg[&\frac{\alpha_s(m_b)}{9\pi}\left\{C_2^q(m_b)-\frac{7}{6}\frac{\alpha}{\alpha_s(m_b)}(3C_1^q(m_b)+C_2^q(m_b))\right\}\mathcal{I}(m_q,k,m_b)\\&-\frac{\alpha_s(m_b)}{9\pi}\left\{4C_8^G(m_b)-7\frac{\alpha}{\alpha_s(m_b)}C_7^{\gamma}(m_b)\right\}+\frac{4}{3}(C_3(m_b)+C_4(m_b))\\&+C_5(m_b)+\frac{1}{3}C_6(m_b)-\frac{1}{2}C_7(m_b)-\frac{1}{6}C_8(m_b)-\frac{2}{3}(C_9(m_b)+C_{10}(m_b))\Bigg]X^{(B^-K^-,\phi)}\\&+(C_i\to C_i'),\end{aligned}\tag{15}$$

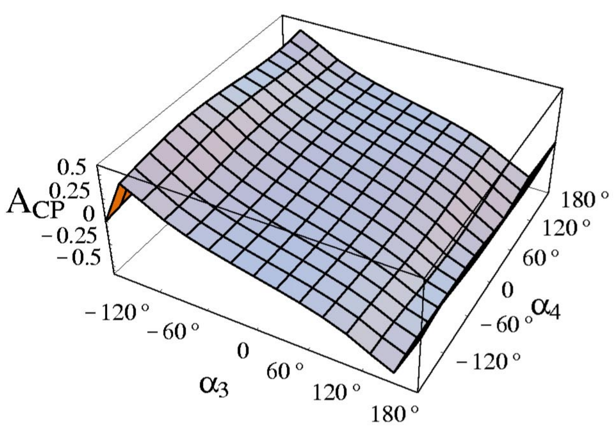
where $X^{(B^-K^-,\phi)}\equiv\langle\phi|\bar{s}\gamma_\mu s|0\rangle\langle K^-|\bar{s}\gamma^\mu b|B^-\rangle$. The amplitude $\mathcal{A}(B^+\to\phi K^+)$ is simply obtained from $\mathcal{A}(B^-\to\phi K^-)$ by replacing $\lambda_q^{LL}\to\lambda_q^{LL*}$ and $C_i^{(\prime)}\to C_i^{(\prime)*}$. In the SM, nonzero *CP* asymmetry arises from the superposition of the *CP*-odd phase γ in V_{ub}^L and the different *CP*-even phases arising from the function $\mathcal{I}(m_q,k,m_b)$ due to the mass difference between c and u quark. The resulting *CP* asymmetry is known to be very small $\sim O(10^{-2})$ [20,24] because the magnitude of the absorptive part is much smaller than that of the dispersive part. Using the numbers in Eq. (9), $m_c=1.3$ GeV, and $\mathrm{Arg}[V_{ub}^L]=-59°$, we can estimate the SM-value of *CP* asymmetry:

$$A_{CP}^{SM}(B^{\pm}\to\phi K^{\pm})\simeq7.3\times10^{-3}.\tag{16}$$

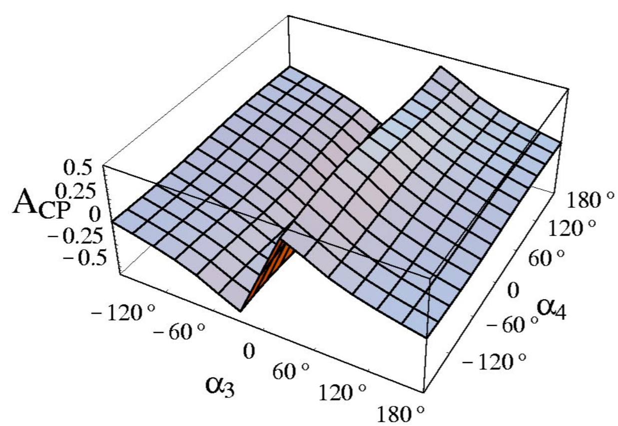
If the model has manifest left-right symmetry, the W_R mass has a stringent bound $M_{W_R}\geqslant1.6$ TeV [25], and its contribution to the decay amplitude is very small so that *CP* asymmetry in the manifest LRM should be very small as well. Since this value is small and our purpose is to estimate the possible large right-handed current contribution, we take a limit $\mathcal{I}(m_c,k,\mu)=\mathcal{I}(m_u,k,\mu)$ in order to get around the uncertainty of V_{ub}^L obtained under the SM framework and clearly see the right-handed current contribution. Then we can express $\mathcal{A}(B^-\to\phi K^-)$ in terms of new parameters ζ_g, ξ_g, and θ_R for two types of V^R in Eq. (4) in the LRM using the unitarity relation $\Sigma_{q=u,c,t}\lambda_q=0$ and the numbers in Eq. (9) again as follows:

$$\begin{aligned}\mathcal{A}(B^-\to\phi K^-)_I\simeq-\frac{G_F}{\sqrt{2}}\{&-2.87e^{i\varphi_1}+23.1e^{i\varphi_2}\zeta_g c_R s_R\\&\times e^{i(\alpha_4-\alpha_3)}+10.1\xi_g(c_R e^{i\alpha_4}-25 s_R e^{i\alpha_3})\}\\&\times10^{-3}X^{(B^-K^-,\phi)},\end{aligned}\tag{17}$$

$$\begin{aligned}\mathcal{A}(B^-\to\phi K^-)_{II}\simeq-\frac{G_F}{\sqrt{2}}\{&-2.87e^{i\varphi_1}+10.1\xi_g c_R e^{i\alpha_4}\}\\&\times10^{-3}X^{(B^-K^-,\phi)},\end{aligned}$$



(a) $B^{\pm}\to\phi K^{\pm}$ decays



(b) $B^{\pm}\to\phi K^{*\pm}$ decays

FIG. 2. Behavior of A_{CP} as $\alpha_{3,4}$ are varied in the case of V_I^R.

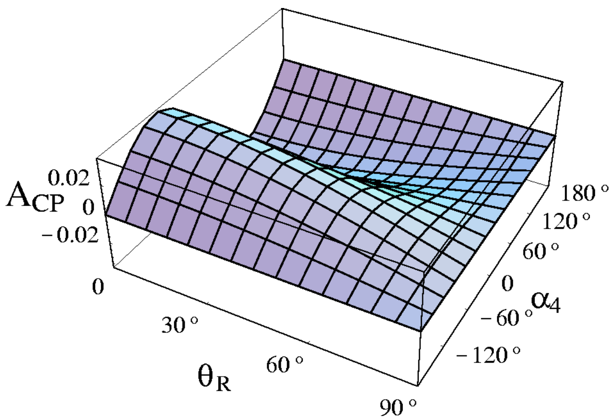


FIG. 3. Behavior of $A_{CP}(B^\pm \to \phi K^{(*)\pm})$ as θ_R and α_4 are varied in the case of V_{II}^R.

where $(\varphi_1, \varphi_2) = (-14.9°, -53.1°)$ are *CP*-even phases. As stated earlier, one can clearly see here that the ζ_g term coming from the coefficients $C'_{1,2}$ is not negligible in case of V_I^R. Likewise, the transition amplitude in $B^- \to \phi K^{*-}$ decays can be easily obtained by using $\langle O_i \rangle = -\langle O'_i \rangle$ because K^{*-} is a vector particle:

$$\mathcal{A}(B^- \to \phi K^{*-})_I \simeq -\frac{G_F}{\sqrt{2}}\{2.87 e^{i\varphi_1} + 23.1 e^{i\varphi_2}\zeta_g c_R s_R \times e^{i(\alpha_4-\alpha_3)} + 10.1\xi_g(-c_R e^{i\alpha_4} - 25 s_R e^{i\alpha_3})\} 10^{-3} X^{(B^- K^{*-},\phi)}, \quad (18)$$

$$\mathcal{A}(B^- \to \phi K^{*-})_{II} \simeq -\frac{G_F}{\sqrt{2}}\{2.87 e^{i\varphi_1} - 10.1\xi_g c_R e^{i\alpha_4}\} \times 10^{-3} X^{(B^- K^{*-},\phi)},$$

where $X^{(B^- K^{*-},\phi)} \equiv \langle \phi | \bar{s}\gamma_\mu s|0\rangle \langle K^{*-}|\bar{s}\gamma^\mu \gamma_5 b|B^-\rangle$. Although the *CP* asymmetry in $B^- \to \phi K^-$ decays should be the same as that in $B^- \to \phi K^{*-}$ decays in the SM, they can be different in the LRM so that the measured difference of *CP* asymmetries between them may give the size of the NP effects.

The current data on the *CP* asymmetries in $B^- \to \phi K^-$ and $B^- \to \phi K^{*-}$ decays are [26]

$$A_{CP}^{\mathrm{expt}}(B^\pm \to \phi K^\pm) = 0.05 \pm 0.20 \pm 0.03,$$

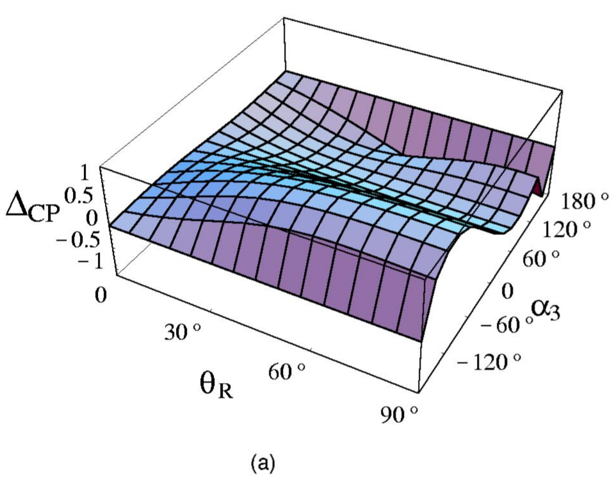
$$A_{CP}^{\mathrm{expt}}(B^\pm \to \phi K^{*\pm}) = 0.43^{+0.36}_{-0.30} \pm 0.06. \quad (19)$$

The SM value in Eq. (16) lies in the range of $A_{CP}^{\mathrm{expt}}(B^\pm \to \phi K^\pm)$, but a little off the range of $A_{CP}^{\mathrm{expt}}(B^\pm \to \phi K^{*\pm})$. In order to explicitly compare these values with the theoretical estimates in the LRM, we first plot $A_{CP}(B^\pm \to \phi K^\pm)$ and $A_{CP}(B^\pm \to \phi K^{*\pm})$ in the case of V_I^R in Fig. 2 for the typical values $\zeta_g = 0.01$, $\xi_g = 0.008$, and $\theta_R = 70°$ as $\alpha_{3,4}$ are varied. In the figure, *CP* asymmetry is drastically changing by varying α_3, and this behavior holds for other values of ζ_g, ξ_g, and θ_R. For the given inputs, $A_{CP}(B^\pm \to \phi K^\pm)$ and $A_{CP}(B^\pm \to \phi K^{*\pm})$ can be different by about 0.5. In the case of V_{II}^R, one can see from Eqs. (18), (19) that $A_{CP}(B^\pm \to \phi K^\pm) = A_{CP}(B^\pm \to \phi K^{*\pm})$ because it has no dependence of ζ_g and α_3 unlike the previous case. In Fig. 3, we fix $\xi_g = 0.01$, and evaluate *CP* asymmetry by varying θ_R and α_4. It shows that *CP* asymmetry is very small with a small parameter ξ_g. Therefore, if we observe large *CP* asymmetry or any difference between $A_{CP}(B^\pm \to \phi K^\pm)$ and $A_{CP}(B^\pm \to \phi K^{*\pm})$, the second type of mass mixing matrix V_{II}^R is disfavored.

B. Neutral *B* meson decays

In the case of the neutral *B* meson decays into *CP* self-conjugate final states *f*, mixing induced *CP* asymmetry can be expressed by the parametrization invariant quantity λ defined by [1]

$$\lambda \equiv \eta_f \left(\frac{q}{p}\right)_B \frac{\mathcal{A}(\overline{B^0} \to \overline{f})}{\mathcal{A}(B^0 \to f)}, \quad \left(\frac{q}{p}\right)_B \simeq \frac{M_{12}^*}{|M_{12}|}, \quad (20)$$



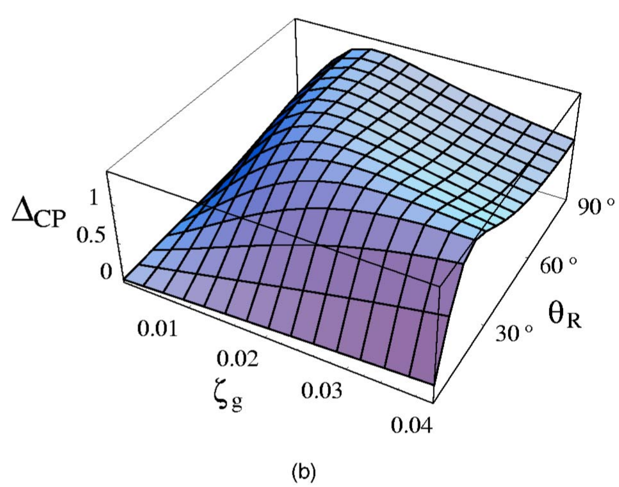


FIG. 4. Behavior of the *CP* asymmetry difference Δ_{CP} between $B \to J/\psi K_S$ and $B \to \phi K_S$ decays in the case of V_I^R.

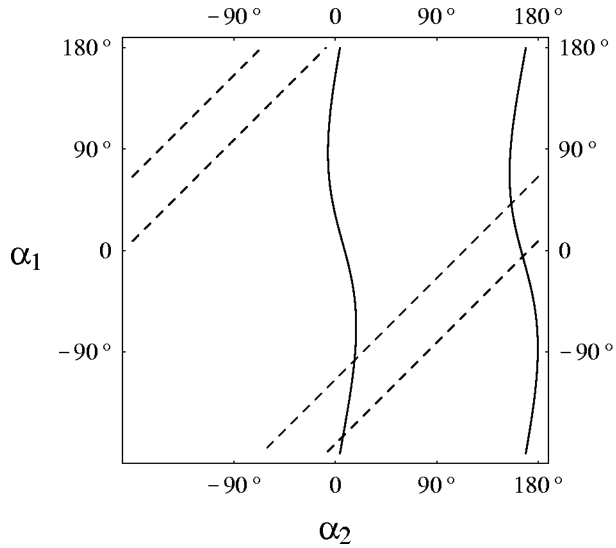


FIG. 5. Contour plot corresponding to $\mathrm{Im}\,\lambda(B\to J/\psi K_S)=0.73$ (solid line) and $\mathrm{Im}\,\lambda(B\to\phi K_S)=-0.39$ (dashed line) for $\sin 2\beta=0.64$ in the case of V_I^R.

where $\eta_f=1(-1)$ for a CP-even(odd) final state f and M_{12} is the dispersive part of the $B\bar{B}$ mixing matrix element. The CP angle β mentioned earlier is simply the imaginary part of λ in $B\to J/\psi K_S$ decays in the SM:

$$\sin 2\beta=\mathrm{Im}\,\lambda(B\to J/\psi K_S)\simeq\mathrm{Im}\,\lambda(B\to\phi K_S). \tag{21}$$

In the general LRM, M_{12} can be written as

$$M_{12}=M_{12}^{SM}+M_{12}^{LR}=M_{12}^{SM}\{1+r_{LR}\}, \tag{22}$$

where

$$r_{LR}\equiv\frac{M_{12}^{LR}}{M_{12}^{SM}}=\frac{\langle\overline{B^0}|H_{eff}^{LR}|B^0\rangle}{\langle\overline{B^0}|H_{eff}^{SM}|B^0\rangle}, \tag{23}$$

with the effective Hamiltonian $H_{eff}^{B\bar{B}}=H_{eff}^{SM}+H_{eff}^{LR}$ in the $B\bar{B}$ system. Considering the two types of the quark mixing matrices in Eq. (4), the effective Hamiltonians in the $B\bar{B}$ system are given by

$$H_{eff}^{SM}=\frac{G_F^2M_W^2}{4\pi^2}(\lambda_t^{LL})^2S(x_t^2)(\overline{d_L}\gamma_\mu b_L)^2, \tag{24}$$

$$\begin{aligned}H_{eff}^{LR}=\frac{G_F^2M_W^2}{2\pi^2}[&\{\lambda_c^{LR}\lambda_t^{RL}x_cx_t\zeta_gA_1(x_t^2,\zeta)\\&+\lambda_t^{LR}\lambda_t^{RL}x_t^2\zeta_gA_2(x_t^2,\zeta)\}(\overline{d_L}b_R)(\overline{d_R}b_L)\\&+\lambda_t^{LL}\lambda_t^{RL}x_b\xi_g\{x_t^3A_3(x_t^2)(\overline{d_L}\gamma_\mu b_L)(\overline{d_R}\gamma_\mu b_R)\\&+x_tA_4(x_t^2)(\overline{d_L}b_R)(\overline{d_R}b_L)\}],\end{aligned} \tag{25}$$

where $S(x)$ is the usual Inami-Lim function and A_i can be found in Ref. [5]. If we consider QCD effect in $B\bar{B}$ mixing, the correction factors should be included in the functions S and A_i. However, there are many uncertainties such as hadronic matrix elements and new parameters in the LRM to prevent us from the precision analysis at this stage, and the QCD corrections to $B\bar{B}$ mixing are not big enough to change our numerical estimate. Therefore we will ignore the QCD corrections to $B\bar{B}$ mixing for simplicity. In the case of V_I^R, there is no significant contribution of H_{eff}^{LR} to $B\bar{B}$ mixing, so that $M_{12}=M_{12}^{SM}$ because $\lambda_t^{RL}\simeq 0$. In the case of V_{II}^R, using $m_c=1.3$ GeV, $m_b=4.4$ GeV, $m_t=170$ GeV, and $|V_{cd}^L|\approx 0.224$, and adopting the parametrization of the hadronic matrix elements of the operators given in Ref. [5], one can express r_{LR} in terms of the mixing angle and phases in Eq. (4) as

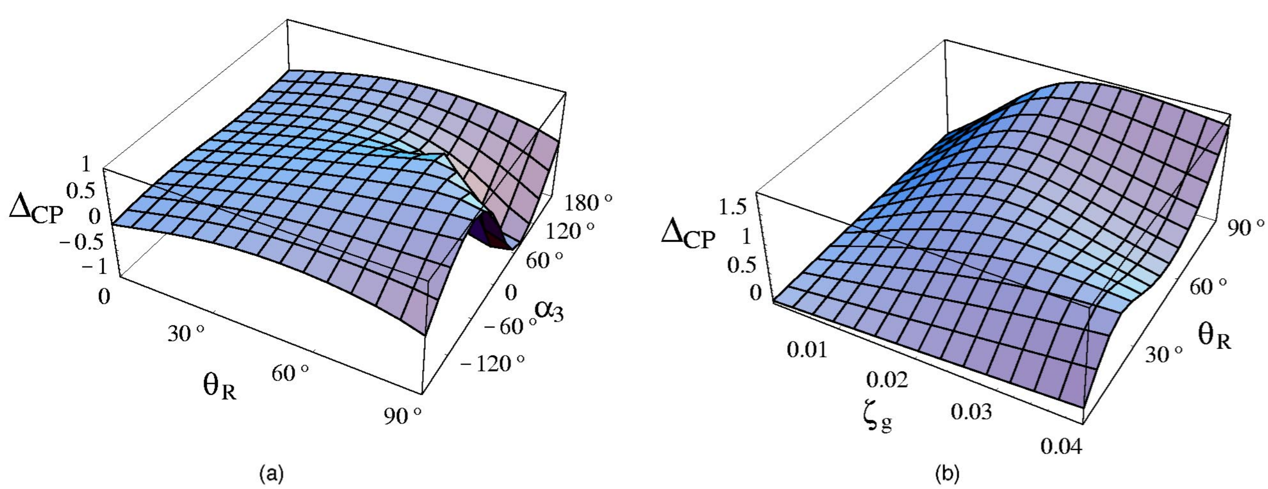


FIG. 6. Behavior of the CP asymmetry difference Δ_{CP} between $B\to J/\psi K_S$ and $B\to\phi K_S$ decays in the case of V_{II}^R.

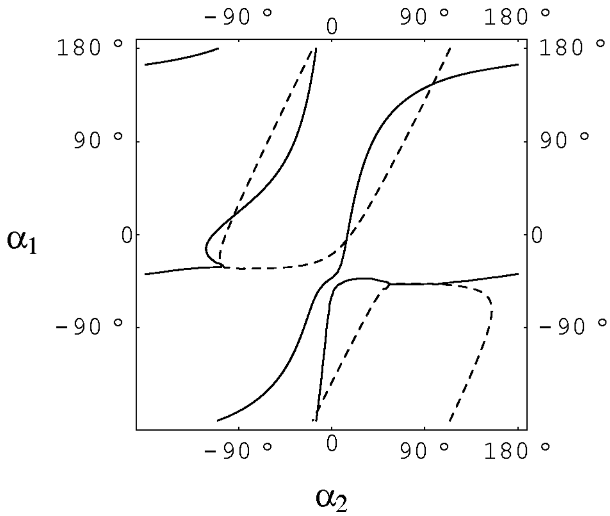


FIG. 7. Contour plot corresponding to $\mathrm{Im}\,\lambda(B\to J/\psi K_S)=0.73$ (solid line) and $\mathrm{Im}\,\lambda(B\to\phi K_S)=-0.39$ (dashed line) for $\sin 2\beta=0.64$ in the case of V_{II}^R.

$$r_{LR}\approx l\Biggl\{17.3l\left(\frac{1-\zeta_g-(4.92-19.7\zeta_g)\ln(1/\zeta_g)}{1-5.47\zeta_g}\right)\zeta_g s_R^2 e^{i\delta_1} -796\left(\frac{1-5.02\zeta_g-(0.498-1.99\zeta_g)\ln(1/\zeta_g)}{1-9.94\zeta_g+28.9\zeta_g^2}\right) \times\zeta_g s_R c_R e^{i\delta_2}-8.93\xi_g s_R e^{i\delta_3}\Biggr\}, \tag{26}$$

where $l=0.008/|V_{td}^L|$, $\delta_1=-2\beta+\alpha_2-\alpha_3$, $\delta_2=-\beta-\alpha_3+\alpha_4$, $\delta_3=-\beta-\alpha_3$. Since $B\to J/\psi K_S$ decay is governed by the tree-level amplitude, the transition amplitude is given by

$$\mathcal{A}(B\to J/\psi K_S)_I\simeq\frac{G_F}{\sqrt{2}}\lambda_c^{LL}\{1+25(c_R s_R\zeta_g e^{-i(\alpha_2-\alpha_1)} -2s_R\xi_g e^{-i\alpha_2})\}X^{(BK_S,J/\psi)},$$

$$\mathcal{A}(B\to J/\psi K_S)_{II}\simeq\frac{G_F}{\sqrt{2}}\lambda_c^{LL}\{1-50 s_R\xi_g e^{-i\alpha_2}\}X^{(BK_S,J/\psi)}, \tag{27}$$

where $X^{(BK_S,J/\psi)}\equiv\langle J/\psi|\bar{c}\gamma_\mu c|0\rangle\langle K_S|\bar{s}\gamma^\mu b|B^\circ\rangle$, and we ignored the $K\bar{K}$ mixing. The transition amplitude in $B\to\phi K_S$ decays can be simply obtained from Eq. (18) by replacing the hadronic matrix element $X^{(B^-K^-,\phi)}\to X^{(BK_S,\phi)}$.

For illustration of the possible effect of the new interaction on the mixing induced *CP* asymmetry, we assume that $\beta=20°$ and $l=1$, and show that the region of parameters α_i where $\mathrm{Im}\,\lambda(B\to J/\psi K_S)\simeq0.73$ and $\mathrm{Im}\,\lambda(B\to\phi K_S)\simeq-0.39$ since $|\lambda|\approx1$. To do so, we need to find an appropriate set of parameters ζ_g, ξ_g, and θ_R yielding a large difference $\Delta_{CP}\equiv\mathrm{Im}\,\lambda(B\to J/\psi K_S)-\mathrm{Im}\,\lambda(B\to\phi K_S)$. First, we evaluate Δ_{CP} in the case of V_I^R for $\zeta_g=\xi_g=0.01$, $\alpha_{1,2}\alpha1,2=0$ by varying θ_R and α_3 in Fig. 4(a). In the figure, Δ_{CP} becomes maximal near $\alpha_3\sim-120°$ and increases as θ_R increases, and this behavior holds for other values of fixed parameters. Since we assumed that Δ_{CP} is larger than 1, we fix $\alpha_3=-120°$, and evaluate Δ_{CP} in Fig. 4(b) for $\alpha_{1,2}=0$ and $\xi_g=\zeta_g$ by varying θ_R and ζ_g. One can see from the figure that Δ_{CP} approaches 1 for $\zeta_g\gtrsim0.01$ and $\theta_R\gtrsim10°$, and its variation is small. After repeating this analysis, we get a probable set of parameter values $\zeta_g=0.01$, $\xi_g=0.008$, $\theta_R=70°$, and $\alpha_3=-120°$. Using these values, we plot the contours corresponding to $\mathrm{Im}\,\lambda(B\to J/\psi K_S)=0.73$ and $\mathrm{Im}\,\lambda(B\to\phi K_S)=-0.39$ in the parameter space of $\alpha_{1,2}$ in Fig. 5. Therefore, as a result from the obtained figures, the manifest or pseudomanifest LRM is disfavored under the given assumption. In a similar way to the case of V_I^R, the results of the analysis of the mixing induced *CP* asymmetries in the case of V_{II}^R are represented in Fig. 6 and Fig. 7.

IV. CONCLUSIONS

In this paper, we studied *CP* asymmetries in penguin-induced $b\to s\bar{s}s$ decays in the general LRM. Without imposing manifest or pseudomanifest left-right symmetry, one has two types of mass mixing matrix V^R with which the right-handed current contributions to $B\bar{B}$ mixing and *CP* asymmetry can be sizable even in the decays such as $B^\pm\to\phi K^{(*)\pm}$ decays where the SM contribution to *CP* asymmetry is very small. Using the effective Hamiltonian approach, we evaluate the sizes of the NP contributions to *CP* asymmetries in $B^\pm\to\phi K^{(*)\pm}$ decays, and show that V_I^R is more probable than V_{II}^R if *CP* asymmetries in those decays are large or different from each other. Similar argument can be made in mixing induced *B* decays such as $B\to J/\psi K_S$ and $B\to\phi K_S$ decays. Although SM predicts that the *CP* asymmetry in $B\to J/\psi K_S$ decays should be very close to that in $B\to\phi K_S$ decays, the present experiments show a large discrepancy between them. Based on these preliminary experimental results, we find that the manifest or pseudomanifest LRM is disfavored, and the bounds of the new parameters are restricted as shown in Figs. 4–7. Furthermore, this result may affect the sizes of *CP* asymmetries in other decays. For instance, one can see from Fig. 2 and Fig. 3 that the contributions of the obtained parameter sets from Fig. 5 and Fig. 7 under the given assumption reduces the size of *CP* asymmetries in $B^\pm\to\phi K^{*\pm}$ decays. In this way, *CP* asymmetries in other mixing induced decays such as $B\to\phi K^*$ can be estimated systematically, and all of these analysis of possible NP contributions can be tested once the experimental results are confirmed.

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