Right-handed current contributions in $B \to K\pi$ decays

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Abstract

The current measurements of CP asymmetries in $B \to K\pi$ decays are in disagreement with the predictions of the Standard Model. In order to solve this discrepancy, using the effective Hamiltonian approach, we investigate the right-handed current contributions to $B \to K\pi$ decay amplitudes including all possible low-energy operators in the nonmanifest left-right model. We find the allowed region of new physics parameters satisfying the current experimental data, and discuss its implication to other observables such as B_s mixing and the branching fraction for $B \to \tau\nu$ decays.

CP asymmetry measurements in non-leptonic $b \to s$ decays have been receiving considerable attention over the past several years since large discrepancies have been observed in some decay channels between the recent experimental measurements and the Standard Model (SM) predictions. One of the important examples is the direct CP asymmetries in $B \to K\pi$ decays [1]. Current world averages of the CP asymmetries in $B \to K\pi$ decays are given by [2]:

$$\begin{aligned} A_{CP}(B^0 \to K^\pm\pi^\mp) &= -0.098 \pm 0.012, \\ A_{CP}(B^\pm \to K^\pm\pi^0) &= 0.050 \pm 0.025, \\ A_{CP}(B^\pm \to K^0\pi^\pm) &= 0.009 \pm 0.025. \end{aligned} \tag{1}$$

However, the naive factorization assumption predicts: $A_{CP}(B^0 \to K^\pm\pi^\mp) \approx A_{CP}(B^\pm \to K^\pm\pi^0)$ [3], which is inconsistent with the current data in Eq. (1). These decay modes are also studied within the SM in the framework of different factorization approaches such as QCDF [4], PQCD [5], and SCET [6]. In the SM, the sizes and patterns of CP violation in various decay modes are governed by a single complex phase which resides in the Cabibbo-Kobayashi-Maskawa (CKM) matrix, but such large CP violation effects has not been simply explained with this single parameter in any of those factorization methods in various decay modes simultaneously. Therefore, there has been several efforts to resolve this issue beyond the SM with additional CP odd parameters [7]. One of the simplest extensions of the SM corresponding to such a scenario is the nonmanifest ($V^R \neq V^L$) left-right model (LRM) with gauge group $SU(2)_L \times SU(2)_R \times U(1)$ where $V^L(V^R)$ is the left(right)-handed quark mixing matrix [8]. In general, the form of V^R is not restricted, but we take the following form with which the right-handed current contributions to CP violating observables in $B \to K\pi$ decays can be sizable:

$$V^R = \begin{pmatrix} \sim 0 & c_R e^{i\alpha_1} & s_R e^{i\alpha_2} \\ e^{i\omega} & \sim 0 & \sim 0 \\ \sim 0 & -s_R e^{i\alpha_3} & c_R e^{i\alpha_4} \end{pmatrix}, \tag{2}$$

where c_R (s_R) $\equiv \cos\theta_R$ ($\sin\theta_R$) ($0^\circ \leq \theta_R \leq 90^\circ$). Here the matrix elements indicated as ~ 0 may be $\lesssim 10^{-2}$ and unitarity requires $\alpha_1 + \alpha_4 = \alpha_2 + \alpha_3$. Especially, with this form, a large CP-violating phase in B_s mixing observed at Tevatron can be naturally accommodated without fine-tuning of new physics (NP) parameters [9], and the present experimental measurement of the large branching fractions for $B \to \tau\nu$ decays can be explained [10].

Since the LRM has the extended group $SU(2)_R$, there are new parameters such as a right-handed gauge coupling g_R, new charged (neutral) gauge bosons W_R (Z_R), and the $W_L - W_R$ ($Z_L - Z_R$) mixing angle ξ (η). After spontaneous

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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. Similarly, the neutral gauge bosons mix each other [11], but we do not present them here because Z_R contribution to flavor changing B decays is negligible. Although tree-level flavor-changing neutral Higgs bosons with masses M_H enter into our theory due to gauge invariance, we also neglect their contributions by assuming $M_H \gg M_{W'}$ [12]. The 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 [13]. For instance, the lower bound on $M_{W'}$ can be obtained from global analysis of muon decay measurements without imposing discrete left-right symmetry [14]:

$$\zeta_g < 0.031 \qquad \text{or} \qquad M_{W'} > (g_R/g_L) \times 460\ \text{GeV}. \tag{3}$$

where $\zeta_g \equiv g_R^2 M_W^2 / g_L^2 M_{W'}^2$. In general, $\zeta_g \geq \xi_g \equiv (g_R/g_L)\xi$ for ordinary Higgs representations [13, 15].

In order to include QCD effects systematically, we start from the following low energy effective Hamiltonian describing $\Delta B = 1$ and $\Delta S = 1$ transition as done similarly in Ref. [16]:

$$\mathcal{H}_{eff} = \frac{G_F}{\sqrt{2}} \left[\sum_{\substack{i=1,2,11,12 \\ q=u,c}} \lambda_q^{LL} C_i^q O_i^q - \lambda_t^{LL} \left(\sum_{i=3}^{10} C_i O_i + C_7^\gamma O_7^\gamma + C_8^G O_8^G \right) \right] + (C_i O_i \to C_i' O_i'), \tag{4}$$

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]. In addition to those SM operators, 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 $O'_{1,2}$ as follows:

$$O_{11}^q = \left(\bar{s}_\alpha q_\beta\right)_{\text{V-A}} \left(\bar{q}_\beta b_\alpha\right)_{\text{V+A}}, \qquad O_{12}^q = \left(\bar{s}_\alpha q_\alpha\right)_{\text{V-A}} \left(\bar{q}_\beta b_\beta\right)_{\text{V+A}}, \tag{5}$$

where $(V \pm A)$ refers to the Lorentz structure $\gamma_\mu(1 \pm \gamma_5)$. These new operators may play an important role in tree-level dominated b decays in the general LRM.

The low-energy effects of the full theory at an arbitrary low energy scale μ can then be described by the linear combination of the given operators and the corresponding Wilson coefficients (WCs) $C_i(\mu)$. In order to calculate $C_i(\mu)$, we first calculate them at $\mu = M_W$ scale. After performing a straightforward matching computation, we find the WCs including the electromagnetic penguin contributions at W scale neglecting the u-quark mass:

$$\begin{aligned}
C_2^q(M_W) &= 1, \quad C_2^{q\prime}(M_W) = \zeta_g \lambda_q^{RR}/\lambda_q^{LL} \ (q = u, c), \\
C_3(M_W) &= \frac{\alpha}{6\pi} \frac{1}{\sin^2\theta_W} \left[2B(x_t) + C(x_t)\right], \\
C_7(M_W) &= \frac{\alpha}{6\pi} \left[4C(x_t) + D(x_t)\right], \\
C_9(M_W) &= \frac{\alpha}{6\pi} \left[4C(x_t) + D(x_t) + \frac{1}{\sin^2\theta_W}(10B(x_t) - 4C(x_t))\right], \\
C_7^\gamma(M_W) &= F(x_t) + \frac{m_t}{m_b} A^{tb} \widetilde{F}(x_t), \quad C_7^{\gamma\prime}(M_W) = \frac{m_t}{m_b} A^{ts*} \widetilde{F}(x_t), \\
C_8^G(M_W) &= G(x_t^2) + \frac{m_t}{m_b} A^{tb} \widetilde{G}(x_t), \quad C_8^{G\prime}(M_W) = \frac{m_t}{m_b} A^{ts*} \widetilde{G}(x_t), \\
C_{12}^u(M_W) &= A^{ub}, \quad C_{12}^{u\prime}(M_W) = A^{us*},
\end{aligned} \tag{6}$$

where

$$x_U = \frac{m_U^2}{M_W^2} \ (U = u, c, t), \qquad A^{UD} = \xi_g \frac{V_{UD}^R}{V_{UD}^L} e^{i\alpha_\circ} \ (D = b, s), \tag{7}$$

and $\alpha_\circ$ is a CP phase residing in the vacuum expectation values, which can be absorbed in α_i in Eq. (2) by redefining $\alpha_i + \alpha_\circ \to \alpha_i$. All other coefficients are negligible or vanish. In Eq. (6), the explicit forms of the functions $B(x_t)$, $C(x_t)$, and $D(x_t)$ can be found in Refs. [17, 18], and $F(x_t)$, $\widetilde{F}(x_t)$, $G(x_t)$, and $\widetilde{G}(x_t)$ are given in Ref. [19]. In the above magnetic coefficients, the terms proportional to ξ_g and ζ_g 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 coefficient C_2' is not neglected because $\zeta_g \geq \xi_g$ and there is possible enhancement by the ratio of CKM angles $(\lambda_q^{RR}/\lambda_q^{LL})$ in the nonmanifest LRM. Note that the new coefficient $C_{12}^{u(\prime)}(M_W)$ can be important in some $b \to s$ transitions because the ξ_g suppression can be offset by the ratio V_{uD}^R/V_{uD}^L in Eq. (7).

The coefficients $C_i(\mu)$ at the scale μ below m_b can be obtained by evolving the coefficients $C_i(M_W)$ with the 28×28 anomalous dimension matrix applying the usual renormalization group procedure in the following way:

$$\vec{C}(\mu) = U_4(\mu, m_b)M(m_b)U_5(m_b, M_W)\vec{C}(M_W), \tag{8}$$

where U_f is the evolution matrix for f active flavors and $M(m)$ gives the matching corrections between $\vec{C}_{f-1}(m)$ and $\vec{C}_f(m)$. In the LL approximation $M = 1$ and the evolution matrix $U(m_1, m_2)$ is given by

$$U(m_1, m_2) = V\left[\left(\frac{\alpha_s(m_2)}{\alpha_s(m_1)}\right)^{\vec{\gamma}/2\beta_0} - \frac{\alpha}{2\beta_0}K(m_1, m_2)\right]V^{-1}, \tag{9}$$

where V diagonalizes the transposed of the anomalous dimension matrix γ_s and $\vec{\gamma}$ is the vector containing the eigenvalues of γ_s^T. In the right-hand side of Eq. (9), the first term represents the pure QCD evolution and the second term describes the additional evolution in the presence of the electromagnetic interaction. The leading order formula for the matrix $K(m_1, m_2)$ can be found in Refs. [17, 18]. Since the strong interaction preserves chirality, the 28×28 anomalous dimensional matrix decomposes into two identical 14×14 blocks. The SM 12×12 submatrix describing the mixing among $O_1 - O_{10}$, O_7^γ, and O_8^G can be found in Ref. [20], and the explicit form of the remaining 4×4 matrix describing the mixing among $O_{11,12}$, O_7^γ, and O_8^G, which partially overlaps with the SM 12×12 submatrix, can be found in Ref. [19].

In this letter, unlike the previous analysis in Ref. [16], we set the scale of weak WCs at $\mu = 1.5$ GeV to use the PQCD results for the hadronic matrix elements [21]. For 4 flavors, we have the following numerical values of C_i(1.5 GeV) in LL precision using the standard quark masses:[1]

$$\begin{gathered}
C_1^q = -0.453, \qquad C_1^{q\prime} = C_1^q \zeta_g \lambda_q^{RR}/\lambda_q^{LL}, \\
C_2^q = 1.231, \qquad C_2^{q\prime} = C_2^q \zeta_g \lambda_q^{RR}/\lambda_q^{LL}, \\
C_3 = 0.024, \quad C_4 = -0.046, \quad C_5 = 0.012, \quad C_6 = -0.066, \\
C_7 = 0.014\alpha, \quad C_8 = 0.069\alpha, \quad C_9 = -1.436\alpha, \quad C_{10} = 0.503\alpha, \\
C_7^\gamma = -0.389 - 17.86A^{tb}, \qquad C_7^{\gamma\prime} = -17.86A^{ts*}, \\
C_8^G = -0.177 - 7.858A^{tb}, \qquad C_8^{G\prime} = -7.858A^{ts*}. \\
C_{11}^u = 0.641A^{ub}, \quad C_{12}^u = 0.879A^{ub}, \quad C_{11}^{u\prime} = 0.641A^{us*}, \quad C_{12}^{u\prime} = 0.879A^{us*},
\end{gathered} \tag{10}$$

where subdominant NP terms are neglected. Note that $C_3' - C_{10}'$ are negligible comparing to $C_7^{\gamma\prime}$ and $C_8^{G\prime}$ whereas $C_{1,2}'$ and $C_{11,12}^{(\prime)}$ are not. $C_{1,2}^{(\prime)}$ and $C_{11,12}^{(\prime)}$ can be important especially to the tree-dominated B decays.

Following the procedure of Ref. [16] of including the penguin-type diagrams of the current-current operators $O_{1,2}$ and the tree-level diagrams associated with the magnetic operators O_7^γ and O_8^G, the one-loop matrix elements of $\mathcal{H}_{eff}$ can be rewritten in terms of the tree-level matrix elements of the effective operators:

$$< sq\bar{q}|\mathcal{H}_{eff}|B >= -\frac{G_F}{\sqrt{2}}\lambda_t^{LL}\sum_{i=1}^{12} C_i^{eff} < sq\bar{q}|O_i|B >^{tree} + (C_i^{eff}O_i \to C_i^{eff\prime}O_i'), \tag{11}$$

with the effective WCs

$$C_i^{eff(\prime)} = C_i^{(\prime)} \qquad (i = 1, 2, 8, 10, 11, 12),$$

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

$$
\begin{aligned}
C_3^{eff(\prime)} &= C_3^{(\prime)} - \frac{1}{N_c} C_g^{(\prime)}, \qquad C_4^{eff(\prime)} = C_4^{(\prime)} + C_g^{(\prime)}, \\
C_5^{eff(\prime)} &= C_5^{(\prime)} - \frac{1}{N_c} C_g^{(\prime)}, \qquad C_6^{eff(\prime)} = C_6^{(\prime)} + C_g^{(\prime)}, \\
C_7^{eff(\prime)} &= C_7^{(\prime)} + C_\gamma^{(\prime)}, \qquad C_9^{eff(\prime)} = C_9^{(\prime)} + C_\gamma^{(\prime)},
\end{aligned} \tag{12}
$$

where

$$
\begin{aligned}
C_g^{(\prime)} &= -\frac{\alpha_s}{8\pi}\left[\frac{1}{\lambda_t^{LL}} \sum_{q=u,c} \lambda_q^{LL} C_2^{q(\prime)} \mathcal{I}(m_q, k, m_b) + 2 C_8^{G(\prime)} \frac{m_b^2}{k^2}\right], \\
C_\gamma^{(\prime)} &= -\frac{\alpha}{3\pi}\left[\frac{1}{\lambda_t^{LL}} \sum_{q=u,c} \lambda_q^{LL} (C_1^{q(\prime)} + \frac{1}{N_c} C_2^{q(\prime)}) \mathcal{I}(m_q, k, m_b) + C_7^{\gamma(\prime)} \frac{m_b^2}{k^2}\right],
\end{aligned} \tag{13}
$$

and

$$
\mathcal{I}(m,k,\mu) = 4\int_0^1 dx x(1-x) \ln\Big[\frac{m^2 - k^2 x(1-x)}{\mu^2}\Big]. \tag{14}
$$

and where k is the momentum transferred by the photon or the gluon to the $(q,\bar{q})$ pair. Here k^2 is expected to be typically in the range $m_b^2/4 \le k^2 \le m_b^2/2$ [22], and we will use $k^2 = m_b^2/2$ for our numerical analysis. The expression of the decay amplitudes can be further simplified by combining the effective WCs in the following way:

$$
a_{2i-1}^{(\prime)} = C_{2i-1}^{eff(\prime)} + \frac{1}{N_c} C_{2i}^{eff(\prime)}, \quad a_{2i}^{(\prime)} = C_{2i}^{eff(\prime)} + \frac{1}{N_c} C_{2i-1}^{eff(\prime)} \; (i = 1,2,3), \tag{15}
$$

where the factor $1/N_c$ originates from fierzing the operators $O_i^{(\prime)}$ after adopting the factorization assumption, and N_c is simply equal to the number of colors in the naive factorization approximation based on the vacuum-insertion method [23]. Also, in the PQCD approach, $N_c \approx 3$ as well because of the cancellation between the non-factorizable contributions.

The matrix amplitudes for $B \to K\pi$ decays can then be written in terms of the effective WCs as

$$
\begin{aligned}
\mathcal{A}(\bar{B}^0 \to \pi^0 \bar{K}^0) &= \frac{G_F}{2}\Bigg\{\left[\lambda_u^{LL}(a_1 + \rho_u^\pi a_{11}) + \frac{3}{2}\lambda_t^{LL}(a_7 - a_9)\right] X^{(BK,\pi)} \\
&+ \lambda_t^{LL}\left[a_4 - \frac{1}{2}a_{10} + 2\rho_s^K\left(a_6 - \frac{1}{2}a_8\right)\right] X^{(B\pi,K)} \\
&+ \lambda_t^{LL}\left[a_4 - \frac{1}{2}a_{10} + 2\rho_s^B\left(a_6 - \frac{1}{2}a_8\right)\right] X^{(B,\pi K)}\Bigg\} \\
&+ (a_i \to -a_i'),
\end{aligned} \tag{16}
$$

$$
\begin{aligned}
\mathcal{A}(\bar{B}^0 \to \pi^+ K^-) &= \frac{G_F}{\sqrt{2}}\Bigg\{\Big[\lambda_u^{LL}(a_2 + a_{12}) - \lambda_t^{LL}\Big(a_4 + a_{10} + 2\rho_s^K(a_6 + a_8)\Big)\Big] X^{(B\pi,K)} \\
&- \lambda_t^{LL}\left[a_4 - \frac{1}{2}a_{10} + 2\rho_s^B\left(a_6 - \frac{1}{2}a_8\right)\right] X^{(B,\pi K)}\Bigg\} \\
&+ (a_i \to -a_i'),
\end{aligned} \tag{17}
$$

$$
\begin{aligned}
\mathcal{A}(B^- \to \pi^0 K^-) &= \frac{G_F}{2}\Bigg\{\left[\lambda_u^{LL}(a_1 + \rho_u^\pi a_{11}) + \frac{3}{2}\lambda_t^{LL}(a_7 - a_9)\right] X^{(BK,\pi)} \\
&+ \Big[\lambda_u^{LL}(a_2 + a_{12}) - \lambda_t^{LL}\Big(a_4 + a_{10} + 2\rho_s^K(a_6 + a_8)\Big)\Big] X^{(B\pi,K)} \\
&+ \Big[\lambda_u^{LL}(a_2 - a_{12}) - \lambda_t^{LL}\Big(a_4 + a_{10} + 2\rho_s^B(a_6 + a_8)\Big)\Big] X^{(B,\pi K)}\Bigg\} \\
&+ (a_i \to -a_i'),
\end{aligned} \tag{18}
$$

$$
\mathcal{A}(B^- \to \pi^- \bar{K}^0) = \frac{G_F}{\sqrt{2}}\Bigg\{ - \lambda_t^{LL}\left[a_4 - \frac{1}{2}a_{10} + 2\rho_s^K\left(a_6 - \frac{1}{2}a_8\right)\right] X^{(B\pi,K)}
$$

$$+\left[\lambda_u^{LL}(a_2 - a_{12}) - \lambda_t^{LL}\Big(a_4 + a_{10} + 2\rho_s^B\,(a_6 + a_8)\Big)\right] X^{(B,\pi K)}\Big\}$$
$$+(a_i \to -a_i'), \qquad (19)$$

where

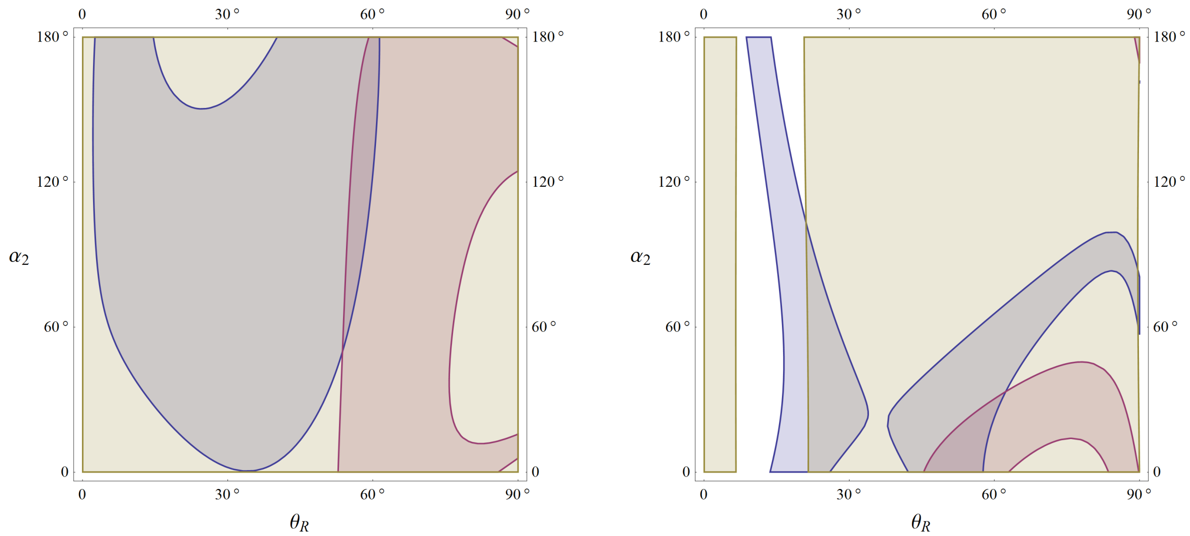
$$\begin{aligned} X^{(BK,\pi)} &= -\sqrt{2} < \pi^0|\bar{u}\gamma^\mu\gamma_5 u|0 >< \bar{K}^0|\bar{s}\gamma_\mu b|\bar{B}^0 > \\ &= i f_\pi F_0^{B\to K}(m_\pi^2)(m_B^2 - m_K^2), \\ X^{(B\pi,K)} &= +\sqrt{2} < \bar{K}^0|\bar{s}\gamma^\mu\gamma_5 d|0 >< \pi^0|\bar{d}\gamma_\mu b|\bar{B}^0 > \\ &= i f_K F_0^{B\to\pi}(m_K^2)(m_B^2 - m_\pi^2), \\ X^{(B,\pi K)} &= +\sqrt{2} < \pi^0\bar{K}^0|\bar{s}\gamma^\mu d|0 >< 0|\bar{d}\gamma_\mu\gamma_5 b|\bar{B}^0 > \\ &= i f_B F_0^{\pi K}(m_B^2)(m_K^2 - m_\pi^2), \end{aligned} \qquad (20)$$

and where

$$\rho_q^H \equiv \frac{m_H^2}{m_b m_q}\ (H = \pi, K, B,\ q = u, s). \qquad (21)$$

Note from Eq. (20) that the form factors $F_0^{B\to K}$ and $F_0^{B\to\pi}$ can be determined by relevant semileptonic B decays, but $F_0^{\pi K}$ is not. Due to the significant hardronic uncertainties in the factorization approximation of the matrix amplitudes, it is very difficult to separately determine the size of NP contributions. Therefore, in this letter, instead of performing a complete analysis by varying all relevant independent NP parameters ($\zeta_g, \xi_g, \theta_R, \alpha_{2,3,4}$) in this model, we fix ξ_g and $\alpha_{1,3,4}$ for simple illustration of NP effects. Also, for numerical analysis, we use the following values of form factors obtained from PQCD calculation [24]:

$$F_0^{B\to K}(m_\pi^2) = 0.37, \quad F_0^{B\to\pi}(m_K^2) = 0.24, \quad F_0^{\pi K}(m_B^2) = (0.39 + 8.16i)10^{-4}. \qquad (22)$$

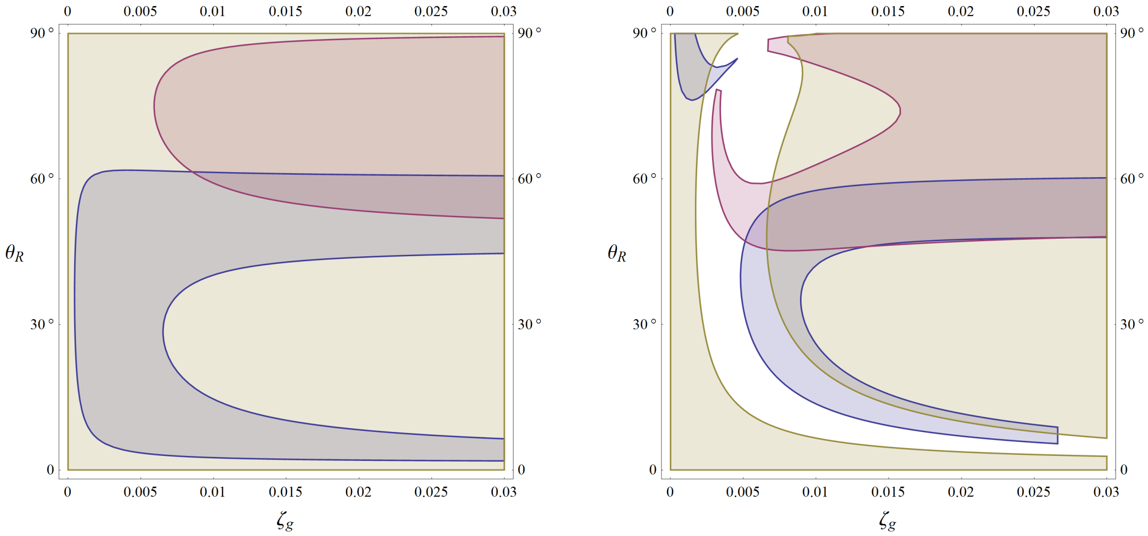


(a) $\alpha_1 = \alpha_3 = 0, \alpha_2 = \alpha_4$

(b) $\alpha_1 = \alpha_2, \alpha_3 = \alpha_4 = \pi$

Figure 1: Allowed regions for α_2 and θ_R at 2 σ level for $M_{W'} = (g_R/g_L) \times 800$ GeV. The blue, red, and green regions are allowed by the current measurements of $A_{CP}(B^0 \to K^\pm\pi^\mp)$, $A_{CP}(B^\pm \to K^\pm\pi^0)$, and $A_{CP}(B^\pm \to K^0\pi^\pm)$, respectively.

Using the present experimental bounds of the CP asymmtries in Eq. (1), we first plot the allowed region of α_2 and θ_R at 2σ level for $\xi_g = \zeta_g = 0.01$ and $\alpha_1 = \alpha_3 = 0$ ($\alpha_3 = \alpha_4 = \pi$) in Fig. 1(a) (Fig. 1(b)). With the chosen NP parameters, the branching fraction of each decay mode in Eq. (16) agrees with the present experimental measurement as well. From the overlapped allowed regions of the figures, one can see that relatively large values of



(a) $\alpha_1 = \alpha_3 = 0, \alpha_2 = \alpha_4 = \pi$

(b) $\alpha_1 = \alpha_2 = 0, \alpha_3 = \alpha_4 = \pi$

Figure 2: Allowed regions for θ_R and ζ_g at 2 σ level. The blue, red, and green regions are allowed by the current measurements of $A_{CP}(B^0 \to K^\pm\pi^\mp)$, $A_{CP}(B^\pm \to K^\pm\pi^0)$, and $A_{CP}(B^\pm \to K^0\pi^\pm)$, respectively.

θ_R are preferred in both cases similarly. This also indicates that manifest ($V^R = V^L$) or pseudomanifest ($V^R = V^{L*}K$) LRM is disfavoured in this case, where K is a diagonal phase matrix. In order to clearly see W' mass dependence, we plot the allowed region of θ_R and ζ_g at 2σ level for $\alpha_2 = \pi$ ($\alpha_2 = 0$) in Fig. 2(a) (Fig. 2(b)). With the given parameter sets, the figures show that the lower bound of ζ_g is approximately in the range of 0.007 ~ 0.01 which corresponds to the upper bound of W' mass which is approximately in the range of 800 GeV ~ 1 TeV. If it happens that the mass of W' is much larger than the obtained upper bound, the right-handed contributions are not big enough to explain the present measurements of the CP asymmetries in $B \to K\pi$ decays.

In summary, we studied the right-handed current contributions to the CP asymmetries in $B \to K\pi$ decays in the nonmanifest left-right model. Without imposing manifest or pseudomanifest left-right symmetry, we parametrized V^R as shown in Eq. (2), and showed that the CP asymmetries are sensitive to the phases and angles in V^R as well as to the mass of W'. With the given phases, relatively large θ_R is preferred as shown in Fig. 1. This could give simultaneous explanations to the large branching fractions of $B \to \tau\nu$ transitions due to a large fraction V^R_{ub}/V^L_{ub} [10] , and also to the large CP violating like-sign dimuon charge asymmetry in semileptonic B decays [25] due to a large CP-violating phase in B_s mixing [9]. Also, with the give parameter sets, Fig. 2 shows that it is favourable that the mass of W' is lighter than around 1 TeV if $g_R = g_L$ in order to incorporate the current experimental measurements. In this way, CP asymmetries in other nonleptonic B decays such as $B \to K\rho$ and $B \to K^*\pi$ can be estimated systematically, and all of these analysis of possible NP contributions can be tested once future experimental progress can further improve the bounds.

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