$K_1(1270) - K_1(1400)$ Mixing Angle and New-Physics Effects in $B \to K_1\ell^+\ell^-$ Decays

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Abstract

We study semileptonic $B$ meson decays $B \to K_1(1270)\ell^+\ell^-$ and $K_1(1400)\ell^+\ell^-$ ($\ell \equiv e, \mu, \tau$), where the strange $P$-wave mesons, $K_1(1270)$ and $K_1(1400)$, are the mixtures of the $K_{1A}$ and $K_{1B}$, which are the $1^3P_1$ and $1^1P_1$ states, respectively. We show that the ratio $R_\ell \equiv B(B \to K_1(1400)\ell^+\ell^-)/B(B \to K_1(1270)\ell^+\ell^-)$, insensitive to new-physics parameters, is suitable for determining the $K_1(1270) - K_1(1400)$ mixing angle, $\theta_{K_1}$. The forward-backward asymmetry shows a weak $\theta_{K_1}$-dependence for $B \to K_1(1270)\mu^+\mu^-$, but relatively strong for $B \to K_1(1400)\mu^+\mu^-$. We investigate model-independent new-physics corrections to operators relevant to the $b \to s\ell^+\ell^-$ electroweak-penguin and weak-box diagrams. Furthermore, for the $B \to K_1(1270)\mu^+\mu^-$ decay the position of the forward-backward asymmetry zero, which is almost independent of the value of $\theta_{K_1}$, can be dramatically changed under variation of new-physics parameters.

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TABLE I: Experimental status of branching fractions (in units of $10^{-6}$) for the decays $B \to K^*(892)\gamma$, $K_1(1270)\gamma$, $K_1(1400)\gamma$ and $B \to K^*(892)\ell^+\ell^−$.

| Mode                  | Exp.(Average) | Ref.   | Mode                  | Exp.(Average) | Ref.   |
|-----------------------|---------------|--------|-----------------------|---------------|--------|
| $K^+(892)\gamma$      | 40.3 ± 2.6    | [2, 3, 4] | $K^0(892)\gamma$     | 40.1 ± 2.0    | [2, 3, 4] |
| $K_1^+(1270)\gamma$   | 43 ± 12       | [5]    | $K_1^0(1270)\gamma$  | < 58          | [5]    |
| $K_1^+(1400)\gamma$   | < 15          | [5]    | $K_1^0(1400)\gamma$  | < 15          | [5]    |
| $K^{*+}(892)e^+e^−$   | 1.23$^{+0.69}_{-0.62}$ | [6, 7] | $K^{*0}(892)e^+e^−$  | 1.11$^{+0.30}_{-0.26}$ | [6, 7] |
| $K^{*+}(892)\mu^+\mu^−$ | 0.78$^{+0.56}_{-0.44}$ | [6, 7] | $K^{*0}(892)\mu^+\mu^−$ | 0.98$^{+0.22}_{-0.21}$ | [6, 7] |

I. INTRODUCTION

$b \to s$ transitions in semileptonic and radiative $B$ meson decays contain rich phenomena relevant to the standard model (SM) and new physics (NP). Semileptonic and radiative $B$ decays involving a vector or axial vector meson have been observed by BABAR, Belle and CLEO (see Table I). The rare flavor-changing neutral-current processes, $b \to s\bar{\ell}\ell$, which proceed through the electroweak-penguin and weak-box diagrams in the SM, may provide a hunting ground to search for the NP effects. For $B \to K^*(892)\ell^+\ell^−$ decays, the forward-backward asymmetry has been measured by BABAR [7] and Belle [8]. Very recently, BABAR [9, 10, 11] has reported the measurements for the longitudinal polarization fraction and forward-backward asymmetry (FBA) of $B \to K^*(892)\ell^+\ell^−$, and for the isospin asymmetry of $B^0 \to K^{*0}(892)\ell^+\ell^−$ and $B^± \to K^{*±}(892)\ell^+\ell^−$ channels. The data may hint at the flipped sign(s) of the Wilson coefficients, e.g., the flipped sign of $c_7^{\text{eff}}$ related to the magnetic dipole operator. To extract the moduli and arguments of the effective Wilson coefficients, it is important to measure various observables in different inclusive and exclusive rare processes. These should be considerably improved at LHCb.

The radiative $B$ decay involving the $K_1(1270)$, the orbitally excited ($P$-wave) state, is recently observed by Belle and other radiative and semileptonic decay modes involving $K_1(1270)$ and $K_1(1400)$ are hopefully expected to be seen soon. Some studies for $B \to K_1\ell^+\ell^−$ have been made recently [12, 13, 14]. Just like $B \to K^*(892)\ell^+\ell^−$ decays [15, 16],...
B → K_1^-ℓ^+ℓ^- decays can offer the good probe to the NP, and are much more sophisticated due to the mixing of the K_{1A} and K_{1B}, which are the 1^3P_1 and 1^1P_1 states, respectively. The physical K_1 mesons are K_1(1270) and K_1(1400), described by

\begin{equation}
\left(\begin{array}{c}
|K_1(1270)\rangle \\
|K_1(1400)\rangle 
\end{array}\right) = M \left(\begin{array}{c}
|K_{1A}\rangle \\
|K_{1B}\rangle 
\end{array}\right), \quad \text{with} \quad M = \left(\begin{array}{cc}
\sin\theta_{K_1} & \cos\theta_{K_1} \\
\cos\theta_{K_1} & -\sin\theta_{K_1}
\end{array}\right).
\end{equation}

The magnitude of \theta_{K_1} was estimated to be |\theta_{K_1}| ≈ 34° ∨ 57° in Ref. [23], 35° ≲ |\theta_{K_1}| ≲ 55° in Ref. [24], and |\theta_{K_1}| = 37° ∨ 58° in Ref. [25]. Nevertheless, the sign of the \theta_{K_1} was not yet determined in these studies. From the study for B → K_1(1270)γ and τ → K_1(1270)ν_τ, we recently obtain [26]

\begin{equation}
\theta_{K_1} = -(34 ± 13)^°,
\end{equation}

where the minus sign of \theta_{K_1} is related to the chosen phase of |K_{1A}\rangle and |K_{1B}\rangle. We adopt the following conventions [26]: f_{K_{1A}} > 0 and f_{K_{1B}}^+ > 0, which are defined by

\begin{equation}
\begin{align*}
\langle 0|\bar{\psi}γ^μγ_5s|K_{1A}(P,λ)\rangle &= -if_{K_{1A}}m_{K_{1A}}ε^{(λ)}_μ, \\
\langle 0|\bar{\psi}σ_{μν}s|K_{1B}(P,λ)\rangle &= if_{K_{1B}}^+ε_{μναβ}ε^{(λ)}_β P^α, \quad \psi ≡ d, u.
\end{align*}
\end{equation}

Within the SM, we have predicted [26]

\begin{align*}
\mathcal{B}(B^- → K_1^- (1270)γ) &= (66^{+50}_{-30}) \times 10^{-6} \left(\frac{m_{b,pole}}{4.90 \text{ GeV}}\right)^2, \\
\mathcal{B}(B^- → K_1^- (1400)γ) &= (6.5^{+12.8}_{-6.3}) \times 10^{-6} \left(\frac{m_{b,pole}}{4.90 \text{ GeV}}\right)^2,
\end{align*}

where m_{b,pole} is the pole mass of the b quark. In the present paper, we study the observables for B → K_1^-ℓ^+ℓ^- decays, including the dilepton mass spectra, decay rates and forward-backward asymmetries. We further show that the mixing angle \theta_{K_1} can be determined from the B → K_1^-ℓ^+ℓ^- decays. In addition to the study of the \theta_{K_1}, we also investigate the model-independent new-physics corrections to the Wilson coefficients c_7^{\text{eff}}, c_9 and c_{10}. The new-physics parameters can be well constrained by the measurement of B → K_1^-ℓ^+ℓ^- forward-backward asymmetry (FBA), where the position of the FBA zero depends very weakly on the value of the \theta_{K_1}. Hence, the position of zero of the differential FBAs depends on the underlying new physics corrections.

This paper is organized as follows. In Sec. II we introduce the effective Hamiltonian and effective operators therein. In Sec. III we give the definitions for B → K_1(1270) and
TABLE II: The Wilson coefficients $c_i(\mu)$ at the scale $\mu = m_{b,\text{pole}}$ in the SM. Here $c_7^{\text{eff}} \equiv c_7 - \frac{1}{3}c_5 - c_6$.

| $c_1$  | $c_2$  | $c_3$  | $c_4$  | $c_5$  | $c_6$  | $c_7^{\text{eff}}$ | $c_9$  | $c_{10}$ |
|--------|--------|--------|--------|--------|--------|---------------------|--------|---------|
| +1.107 | -0.248 | -0.011 | -0.026 | -0.007 | -0.031 | -0.313             | 4.344  | -4.669  |

$B \to K_1(1400)$ form factors. In Sec. [V] we formulate the $B \to K_1\ell^+\ell^-$ decays and discuss determination of the $\theta_{K_1}$ in details. In Sec. [V] we estimate the NP effects in the model-independent way. We summarize the main results in Sec. [VI].

II. THE EFFECTIVE HAMILTONIAN

Neglecting doubly Cabibbo-suppressed contributions, the effective weak Hamiltonian relevant to $b \to s\ell^+\ell^-$ is given by

$$\mathcal{H}_{\text{eff}} = -\frac{4}{\sqrt{2}} G_F V_{tb} V^*_{ts} \sum_{i=1}^{10} c_i(\mu) O_i(\mu),$$

where the Wilson operators $O_i$ for $i = 1, \cdots, 10$ read [27]

$$O_1 = (\bar{s}_\alpha \gamma_\mu L_{c\alpha})(\bar{c}_\beta \gamma^\mu L_{b\beta}), \quad O_2 = (\bar{s}_\alpha \gamma_\mu L_{c\alpha})(\bar{c}_\beta \gamma^\mu L_{b\alpha}),$$

$$O_3 = (\bar{s}_\alpha \gamma_\mu L_{b\alpha}) \sum_q (\bar{q}_\beta \gamma^\mu Q_{q\beta}), \quad O_4 = (\bar{s}_\alpha \gamma_\mu L_{b\alpha}) \sum_q (\bar{q}_\beta \gamma^\mu L_{q\alpha}),$$

$$O_5 = (\bar{s}_\alpha \gamma_\mu L_{b\alpha}) \sum_q (\bar{q}_\beta \gamma^\mu R_{q\beta}), \quad O_6 = (\bar{s}_\alpha \gamma_\mu L_{b\beta}) \sum_q (\bar{q}_\beta \gamma^\mu R_{q\alpha}),$$

$$O_7 = \frac{\alpha_{\text{em}} m_b}{16\pi^2} \bar{s} \sigma_{\mu\nu} R b F_{\mu\nu},$$

$$O_9 = \frac{\alpha_{\text{em}}}{4\pi} (\bar{\ell} \gamma_\mu \ell)(\bar{s} \gamma^\mu L b), \quad O_{10} = \frac{\alpha_{\text{em}}}{4\pi} (\bar{\ell} \gamma_\mu \gamma_5 \ell)(\bar{s} \gamma^\mu L b),$$

with $L = (1 - \gamma_5)/2$, $R = (1 + \gamma_5)/2$, and $\alpha, \beta$ being the $SU(3)$ color indices. The $b \to s\ell^+\ell^-$ decay amplitude is given by

$$\mathcal{M}(b \to s\ell^+\ell^-) = \frac{G_F \alpha_{\text{em}}}{\sqrt{2}} V_{tb} V^*_{ts} \left\{ c_9^{\text{eff}}(\hat{s})[\bar{s}\gamma_\mu L b][\bar{\ell} \gamma^\mu \ell] + c_{10}[\bar{s}\gamma_\mu L b][\bar{\ell} \gamma_5 \ell] - 2\hat{m}_b \ell_7^{\text{eff}} \left[ \bar{s} \sigma_{\mu\nu} \hat{q}^\nu \hat{s} \right] R b \right\},$$

where $\hat{m}_b \equiv \hat{m}_b/m_B$ with $\hat{m}_b \equiv m_b(\hat{m}_b)$ being the $b$ quark mass in the $\overline{\text{MS}}$ scheme, $\hat{s} = q^2/m_B^2$, $q_\mu = (p_+ + p_-)_\mu$ with $p_\pm$ being momenta of the leptons $\ell^\pm$. To next-to-leading order
the running \( \overline{\text{MS}} \) and pole \( b \)-quark masses are related by

\[
\overline{m}_b(\mu) = m_{b,\text{pole}} \left[ 1 - \frac{\alpha_s(\mu) C_F}{4\pi} \left( 4 - 3 \ln \frac{m_{b,\text{pole}}^2}{\mu^2} \right) + \mathcal{O}(\alpha_s^2) \right],
\]

where \( C_F = (N_c^2 - 1)/(2N_c) \) with \( N_c \) being the number of colors. In Eq. (9) we have neglected \( \mathcal{O}(m_s/m_b) \) corrections. \( e_g^\text{eff}(\hat{s}) = c_0 + Y(\hat{s}) \), where \( Y(\hat{s}) = Y(\text{pert})(\hat{s}) + Y(\text{LD}) \) contains both the perturbative part \( Y(\text{pert})(\hat{s}) \) and long-distance part \( Y(\text{LD})(\hat{s}) \). \( Y(\text{pert})(\hat{s}) \) is given by

\[
Y(\text{pert})(\hat{s}) = g(\hat{m}_c, \hat{s}) c_0 - \frac{1}{2} g(1, \hat{s})(4\hat{c}_3 + 4\hat{c}_4 + 3\hat{c}_5 + \hat{c}_6) - \frac{1}{2} g(0, \hat{s})(\hat{c}_3 + 3\hat{c}_4)
+ \frac{2}{9}(3\hat{c}_3 + \hat{c}_4 + 3\hat{c}_5 + \hat{c}_6),
\]

with \( c_0 \equiv \hat{c}_1 + 3\hat{c}_2 + \hat{c}_3 + \hat{c}_4 + 3\hat{c}_5 + \hat{c}_6 \),

and the function \( g(x, y) \) defined in \( [28] \). Here \( \hat{c}_1 - \hat{c}_6 \) are the Wilson coefficients in the leading logarithmic approximation. The relevant Wilson coefficients are collected in Table [II [15, 27]]. \( Y(\text{LD}) \) involves \( B \rightarrow K_1 V(\bar{c}c) \) resonances \( [29, 30, 31] \), where \( V(\bar{c}c) \) are the vector charmonium states. We follow Refs. \( [29, 30] \) and set

\[
Y(\text{LD})(\hat{s}) = -\frac{3\pi}{\alpha_s^2} c_0 \sum_{V=\psi(1s),\ldots} \kappa_V \frac{\hat{m}_V B(V \rightarrow \ell^+\ell^-) \hat{\Gamma}^V_{\text{tot}}}{\hat{s} - \hat{m}_V^2 + i\hat{m}_V \hat{\Gamma}^V_{\text{tot}}},
\]

where \( \hat{\Gamma}^V_{\text{tot}} \equiv \Gamma^V_{\text{tot}}/m_B \) and \( \kappa_V = 2.3 \). The relevant properties of vector charmonium states are summarized in Table [III]

### III. \( B \rightarrow K_1(1270) \) AND \( B \rightarrow K_1(1400) \) FORM FACTORS

The \( \overline{\mathcal{B}}(p_B) \rightarrow \overline{\mathcal{K}}_1(p_{K_1}, \lambda) \) form factors are defined by

\[
\langle \overline{\mathcal{K}}_1(p_{K_1}, \lambda) | \bar{\psi} \gamma_\mu (1 - \gamma_5) b | \overline{\mathcal{B}}(p_B) \rangle
= -i \frac{2}{m_B + m_{K_1}} \epsilon_{\mu\nu\rho\sigma} \varepsilon^{*(\lambda)} \varepsilon^{*(\lambda)} \varepsilon_{\nu \rho \sigma} A^{K_1}(q^2)
- \left[ (m_B + m_{K_1}) \varepsilon^{*(\lambda)} \varepsilon_\mu A^{K_1}(q^2) - (p_B + p_{K_1}) \mu \varepsilon^{*(\lambda)} \cdot p_B \right] \frac{V_0^{K_1}(q^2)}{m_B + m_{K_1}}
+ 2m_{K_1} \varepsilon^{*(\lambda)} \cdot p_B q_\mu \left[ V^{K_1}_{3}(q^2) - V^{K_1}_{0}(q^2) \right],
\]

\[
\langle \overline{\mathcal{K}}_1(p_{K_1}, \lambda) | \bar{\psi} \sigma_{\mu\nu} q^\nu (1 + \gamma_5) b | \overline{\mathcal{B}}(p_B) \rangle
\]
form factors can be parametrized by

\[ \text{TABLE III: Masses, total decay widths and branching fractions of dilepton decays of vector charmonium states}^{32}. \]

| \( V \)          | \( \text{Mass [GeV]} \) | \( \Gamma \text{tot [MeV]} \) | \( \mathcal{B}(V \to \ell^+\ell^-) \) |
|------------------|-------------------------|-------------------------------|-------------------------------------|
| \( J/\Psi(1S) \) | 3.097                   | 0.093                         | \( 5.9 \times 10^{-2} \) for \( \ell = e, \mu \) |
| \( \Psi(2S) \)   | 3.686                   | 0.327                         | \( 7.4 \times 10^{-3} \) for \( \ell = e, \mu \) |
| \( \Psi(3770) \) | 3.772                   | 25.2                          | \( 3.0 \times 10^{-3} \) for \( \ell = \tau \) |
| \( \Psi(4040) \) | 4.040                   | 80                            | \( 9.8 \times 10^{-6} \) for \( \ell = e \) |
| \( \Psi(4160) \) | 4.153                   | 103                           | \( 1.1 \times 10^{-5} \) for \( \ell = e \) |
| \( \Psi(4415) \) | 4.421                   | 62                            | \( 8.1 \times 10^{-6} \) for \( \ell = e \) |

\[ = 2T_1^{K_1}(q^2)\epsilon_{\mu\nu\rho\sigma}\bar{\psi}\gamma_\lambda(p_Bp_{K_1}) \]
\[ -iT_2^{K_1}(q^2)\left[ (m_B^2 - m_{K_1}^2)\epsilon_{\lambda\mu}^{*} - (\epsilon_{\lambda}^{*}\cdot q)(p_B + p_{K_1})_\mu \right] \]
\[ -iT_3^{K_1}(q^2)(\epsilon_{\lambda}\cdot q)\left[ q_\mu - \frac{q^2}{m_B^2 - m_{K_1}^2}(p_{K_1} + p_B)_\mu \right] , \]

where \( q = p_B - p_{K_1} \), \( \gamma_5 \equiv i\gamma^0\gamma^1\gamma^2\gamma^3 \), \( \epsilon^{0123} = -1 \), and \( \bar{\psi} \equiv d, s \). The form factors satisfy the following relations,

\[ V_{3}^{K_1}(0) = V_{0}^{K_1}(0), \quad T_{1}^{K_1}(0) = T_{2}^{K_1}(0), \]
\[ V_{3}^{K_1}(q^2) = \frac{m_B + m_{K_1}}{2m_{K_1}}V_{1}^{K_1}(q^2) - \frac{m_B - m_{K_1}}{2m_{K_1}}V_{2}^{K_1}(q^2) . \]

Because the \( K_1(1270) \) and \( K_1(1400) \) are the mixing states of the \( K_{1A} \) and \( K_{1B} \), the \( \bar{B} \to K_1 \) form factors can be parametrized by

\[ \left( \frac{\langle \bar{K}_1(1270) | \bar{s}\gamma_\mu(1 - \gamma_5)b | \bar{B} \rangle}{\langle \bar{K}_1(1400) | \bar{s}\gamma_\mu(1 - \gamma_5)b | \bar{B} \rangle} \right) = M \left( \frac{\langle \bar{K}_{1A} | \bar{s}\gamma_\mu(1 - \gamma_5)b | \bar{B} \rangle}{\langle \bar{K}_{1B} | \bar{s}\gamma_\mu(1 - \gamma_5)b | \bar{B} \rangle} \right) , \]
\[ \left( \frac{\langle \bar{K}_1(1270) | \bar{s}\sigma_{\mu\nu}q^\nu(1 + \gamma_5)b | \bar{B} \rangle}{\langle \bar{K}_1(1400) | \bar{s}\sigma_{\mu\nu}q^\nu(1 + \gamma_5)b | \bar{B} \rangle} \right) = M \left( \frac{\langle \bar{K}_{1A} | \bar{s}\gamma_\mu(1 + \gamma_5)b | \bar{B} \rangle}{\langle \bar{K}_{1B} | \bar{s}\gamma_\mu(1 + \gamma_5)b | \bar{B} \rangle} \right) , \]
with the mixing matrix $M$ being given in Eq. (11). Thus the form factors $A^{K_1}, V_{1,2}^{K_1}$ and $T_{1,2,3}^{K_1}$ satisfy the following relations:

\[
\begin{align*}
(A^{K_1(1270)}/(m_B + m_{K_1(1270)})) &= M \left( A^{K_1A}/(m_B + m_{K_{1A}}) \right), \\
(A_{K_1(1400)}/(m_B + m_{K_1(1400)})) &= M \left( A_{K_1B}/(m_B + m_{K_{1B}}) \right), \\
(m_B + m_{K_1(1270)})V_{1K_1(1270)} &= M \left( (m_B + m_{K_{1A}})V_{1K_1A} \right), \\
(m_B + m_{K_1(1400)})V_{1K_1(1400)} &= M \left( (m_B + m_{K_{1B}})V_{1K_1B} \right), \\
(V_{2K_1(1270)}/(m_B + m_{K_1(1270)})) &= M \left( V_{2K_1A}/(m_B + m_{K_{1A}}) \right), \\
(V_{2K_1(1400)}/(m_B + m_{K_1(1400)})) &= M \left( V_{2K_1B}/(m_B + m_{K_{1B}}) \right), \\
(T_{1K_1(1270)}) &= M \left( T_{1K_1A} \right), \\
(T_{1K_1(1400)}) &= M \left( T_{1K_1B} \right), \\
(m_B^2 - m_{K_1(1270)}^2)T_{2K_1(1270)} &= M \left( (m_B^2 - m_{K_{1A}}^2)T_{2K_1A} \right), \\
(m_B^2 - m_{K_1(1400)}^2)T_{2K_1(1400)} &= M \left( (m_B^2 - m_{K_{1B}}^2)T_{2K_1B} \right), \\
(T_{3K_1(1270)}) &= M \left( T_{3K_1A} \right), \\
(T_{3K_1(1400)}) &= M \left( T_{3K_1B} \right),
\end{align*}
\]

where we have assumed that $p_{K_1(1270),K_1(1400)}^\mu \simeq p_{K_{1A}}^\mu \simeq p_{K_{1B}}^\mu$. For the numerical analysis, we use the light-cone sum rule (LCSR) results for the form factors [33, 34] which are exhibited in Table IV, where the momentum dependence is parametrized in the three-parameter form:

\[
F(q^2) = \frac{F(0)}{1 - a(q^2/m_B^2) + b(q^2/m_B^2)^2},
\]

IV. $B \rightarrow K_1\ell^+\ell^-$ DECAYS IN THE SM

The decay amplitude for $B \rightarrow K_1\ell^+\ell^-$ which is analogous to the $B \rightarrow K^*(892)\ell^+\ell^-$ decay [15] is given by

\[
\mathcal{M} = \frac{G_F \alpha_{em}}{2\sqrt{2}\pi} V_{ts} V_{tb} m_B \cdot (-i) \left[ T_{\mu}^{(K_1),1} \bar{\ell} \gamma^\mu \ell + T_{\mu}^{(K_1),2} \bar{\ell} \gamma^\mu \gamma_5 \ell \right],
\]

where

\[
T_{\mu}^{(K_1),1} = A_{K_1}(\hat{s}) \epsilon_{\mu\nu\rho\sigma} \varepsilon^\nu_P B_{K_1}^\rho B_{K_1}^\sigma - i B_{K_1}(\hat{s}) \epsilon^*_\mu.
\]
TABLE IV: Form factors for $B \to K_{1A}, K_{1B}$ transitions obtained in the LCSR calculation \cite{33,34} are fitted to the 3-parameter form in Eq. (25).

| $F$    | $F(0)$  | $a$  | $b$  | $F$    | $F(0)$  | $a$  | $b$  |
|--------|---------|------|------|--------|---------|------|------|
| $V_{1}^{BK_{1A}}$ | 0.34 ± 0.07 | 0.635 | 0.211 | $V_{1}^{BK_{1B}}$ | −0.29+0.08−0.05 | 0.729 | 0.074 |
| $V_{2}^{BK_{1A}}$ | 0.41 ± 0.08 | 1.51  | 1.18  | $V_{2}^{BK_{1B}}$ | −0.17+0.05−0.03 | 0.919 | 0.855 |
| $V_{0}^{BK_{1A}}$ | 0.22 ± 0.04 | 2.40  | 1.78  | $V_{0}^{BK_{1B}}$ | −0.45+0.12−0.08 | 1.34  | 0.690 |
| $A^{BK_{1A}}$     | 0.45 ± 0.09 | 1.60  | 0.974 | $A^{BK_{1B}}$     | −0.37+0.10−0.06 | 1.72  | 0.912 |
| $T_{1}^{BK_{1A}}$ | 0.31+0.09−0.05 | 2.01  | 1.50  | $T_{1}^{BK_{1B}}$ | −0.25+0.06−0.07 | 1.59  | 0.790 |
| $T_{2}^{BK_{1A}}$ | 0.31+0.09−0.05 | 0.629 | 0.387 | $T_{2}^{BK_{1B}}$ | −0.25+0.06−0.07 | 0.378 | −0.755 |
| $T_{3}^{BK_{1A}}$ | 0.28+0.08−0.05 | 1.36  | 0.720 | $T_{3}^{BK_{1B}}$ | −0.11 ± 0.02−1.61 | 10.2 |

\begin{align}
+iC_{K_{1}}(\hat{s})(\varepsilon^* \cdot \hat{p}_{B})\hat{\mu} + iD_{K_{1}}(\hat{s})(\varepsilon^* \cdot \hat{p}_{B})\hat{q}_{\mu},
\end{align}

\begin{align}
T_{\mu}^{(K_{1})2} = \mathcal{E}_{K_{1}}(\hat{s})\epsilon_{\mu\nu\rho\sigma}^{*\nu\hat{p}_{B}}\hat{p}_{K_{1}} - i\mathcal{F}_{K_{1}}(\hat{s})\varepsilon_{\mu\nu}^{*} + i\mathcal{G}_{K_{1}}(\hat{s})\varepsilon_{\mu} + i\mathcal{H}_{K_{1}}(\hat{s})\varepsilon_{\mu},
\end{align}

with $\hat{p} = p/m_{B}, \hat{p}_{B} = p_{B}/m_{B}, \hat{q} = q/m_{B}$ and $p = p_{B} + p_{K_{1}}, q = p_{B} - p_{K_{1}} = p_{+} + p_{-}$. Here $A^{K_{1}}(\hat{s}), \cdots, \mathcal{H}^{K_{1}}(\hat{s})$ are defined by

\begin{align}
A^{K_{1}}(\hat{s}) = \frac{2}{1 + \hat{m}_{K_{1}}} \frac{c^{\text{eff}}_{9}(\hat{s})}{\hat{s}} A_{K_{1}}(\hat{s}) + \frac{4\hat{m}_{b}}{\hat{s}} c^{\text{eff}}_{7} T_{1}^{K_{1}}(\hat{s}),
\end{align}

\begin{align}
B^{K_{1}}(\hat{s}) = (1 + \hat{m}_{K_{1}}) \left[ c^{\text{eff}}_{9}(\hat{s}) V_{1}^{K_{1}}(\hat{s}) + \frac{2\hat{m}_{b}}{\hat{s}} (1 - \hat{m}_{K_{1}}) c^{\text{eff}}_{7} T_{2}^{K_{1}}(\hat{s}) \right],
\end{align}

\begin{align}
C^{K_{1}}(\hat{s}) = \frac{1}{1 - \hat{m}_{K_{1}}^{2}} \left[ (1 - \hat{m}_{K_{1}}) c^{\text{eff}}_{9}(\hat{s}) V_{2}^{K_{1}}(\hat{s}) + 2\hat{m}_{b} c^{\text{eff}}_{7} \left( T_{3}^{K_{1}}(\hat{s}) + \frac{1 - \hat{m}_{K_{1}}^{2}}{\hat{s}} T_{2}^{K_{1}}(\hat{s}) \right) \right],
\end{align}

\begin{align}
D^{K_{1}}(\hat{s}) = \frac{1}{\hat{s}} \left[ c^{\text{eff}}_{9}(\hat{s}) \{ (1 + \hat{m}_{K_{1}}) V_{1}^{K_{1}}(\hat{s}) - (1 - \hat{m}_{K_{1}}) V_{2}^{K_{1}}(\hat{s}) - 2\hat{m}_{K_{1}} V_{0}^{K_{1}}(\hat{s}) \} ight.
- 2\hat{m}_{b} c^{\text{eff}}_{7} T_{3}^{K_{1}}(\hat{s}),
\end{align}

\begin{align}
E^{K_{1}}(\hat{s}) = \frac{2}{1 + \hat{m}_{K_{1}}} c_{10} A^{K_{1}}(\hat{s}),
\end{align}

\begin{align}
F^{K_{1}}(\hat{s}) = (1 + \hat{m}_{K_{1}}) c_{10} V_{1}^{K_{1}}(\hat{s}),
\end{align}

\begin{align}
G^{K_{1}}(\hat{s}) = \frac{1}{1 + \hat{m}_{K_{1}}} c_{10} V_{2}^{K_{1}}(\hat{s}),
\end{align}

\begin{align}
H^{K_{1}}(\hat{s}) = \frac{1}{\hat{s}} c_{10} \left[ (1 + \hat{m}_{K_{1}}) V_{1}^{K_{1}}(\hat{s}) - (1 - \hat{m}_{K_{1}}) V_{2}^{K_{1}}(\hat{s}) - 2\hat{m}_{K_{1}} V_{0}^{K_{1}}(\hat{s}) \right],
\end{align}
TABLE V: Input parameters

| B meson mass and lifetimes [32] |  |
|----------------------------------|--|
| $m_B = 5.279$ GeV, $\tau_{B^-} = 1.638 \times 10^{-12}$ sec, $\tau_{B^0} = 1.530 \times 10^{-12}$ sec |  |

| Axial vector meson masses [GeV] |  |
|----------------------------------|--|
| $m_{K_1(1270)} = 1.272$ [32], $m_{K_1}(1400) = 1.403$ [32], $m_{K_{1A}} = 1.31$ [35], $m_{K_{1B}} = 1.34$ [35] |  |

| CKM matrix elements |  |
|---------------------|--|
| $|V_{tb}|V_{ts}^*| = 0.0407^{+0.0009}_{-0.0008}$ [36] |  |

| b quark mass [GeV] |  |
|--------------------|--|
| $m_{b,\text{pole}} = 4.8 \pm 0.2$ |  |

| Gauge couplings and the parameter for the $B$ meson distribution amplitude |  |
|--------------------------------------------------------------------------|--|
| $\alpha_{em} = 1/129$, $\alpha_s(\mu_h) = 0.3$, $\lambda_{B,+}^{-1} = 3 \pm 1$ GeV$^{-1}$ [16] |  |

| $K_1$ decay constants [MeV] [35] |  |
|----------------------------------|--|
| $f^\parallel_{K_{1A}} = 250 \pm 13$, $f^\perp_{K_{1B}} [1$ GeV$] = 190 \pm 10$ |  |

| Gegenbauer moments at the scale 2.2 GeV [35] |  |
|-----------------------------------------------|--|
| $a^2_{K_{1A},\perp} = 0.24^{+0.03}_{-0.21}$, $a_1^{K_{1A},\perp} = -0.84 \pm 0.37$, $a_2^{K_{1A},\perp} = 0.01 \pm 0.15$, $a_1^{K_{1B},\perp} = 0.25^{+0.00}_{-0.26}$, $a_2^{K_{1B},\perp} = -0.02 \pm 0.17$ |  |

with $\hat{m}_{K_1} = m_{K_1}/m_B$. We choose $\hat{s} = \hat{q}^2$ and $\hat{u} \equiv (\hat{p}_B - \hat{p}_-)^2 - (\hat{p}_B - \hat{p}_+)^2$ as the two independent parameters, which are bounded as $4\hat{m}_{K_1}^2 \leq \hat{s} \leq (1-\hat{m}_{K_1})^2$ and $-\hat{u}(\hat{s}) \leq \hat{u} \leq \hat{u}(\hat{s})$, with $\hat{u}(\hat{s}) \equiv \sqrt{\lambda (1 - 4\hat{m}_{K_1}^2/\hat{s})}$, $\lambda \equiv 1 + \hat{m}_{K_1}^2 + \hat{s}^2 - 2\hat{s} - 2\hat{m}_{K_1}^2 (1 + \hat{s})$. We have $\hat{u} = -\hat{u}(\hat{s}) \cos \theta$, where $\theta$ is the angle between the momenta of $\ell^+$ and the $b$ quark in the center-of-mass frame of the lepton pair. We will use the parameters given in Tables IV and V in the numerical analysis.
A. Dilepton mass spectrum

The dilepton invariant mass spectrum of the lepton pair for the $B \to K_1\ell^+\ell^-$ decay is given by

$$
\frac{d\Gamma(B \to K_1\ell^+\ell^-)}{ds} = \frac{G_F^2|V_{tb}V_{ts}|^2}{2\pi m_B^5} \hat{\mathcal{A}}(s) \times \left\{ \frac{|A^{K_1}|^2}{3} \hat{s} \lambda\left(1 + 2 \hat{m}_\ell^2\right) + |\mathcal{E}^{K_1}|^2 \hat{s} \frac{\hat{u}(s)^2}{3} \right\}
$$

$$
+ \frac{1}{4m_{K_1}^2} \left[ |\mathcal{B}^{K_1}|^2 \left( \lambda - \frac{\hat{u}(s)^2}{3} + 8\hat{m}_{K_1}^2(\hat{s} + 2\hat{m}_\ell^2) \right) \right]
$$

$$
+ |\mathcal{F}^{K_1}|^2 \left( \lambda - \frac{\hat{u}(s)^2}{3} + 8\hat{m}_{K_1}^2(\hat{s} - 4\hat{m}_\ell^2) \right) \right] \right]
$$

$$
+ \frac{\lambda}{4m_{K_1}^2} \left[ |\mathcal{C}^{K_1}|^2 \left( \lambda - \frac{\hat{u}(s)^2}{3} \right) + |\mathcal{G}^{K_1}|^2 \left( \lambda - \frac{\hat{u}(s)^2}{3} + 4\hat{m}_\ell^2(2 + 2\hat{m}_{K_1}^2 - \hat{s}) \right) \right]
$$

$$
- \frac{1}{2m_{K_1}^2} \left[ \text{Re} \left( \mathcal{B}^{K_1} \mathcal{C}^{K_1*} \right) \left( \lambda - \frac{\hat{u}(s)^2}{3} \right) (1 - \hat{m}_{K_1}^2 - \hat{s}) \right]
$$

$$
+ \text{Re} \left( \mathcal{F}^{K_1} \mathcal{G}^{K_1*} \right) \left( \left( \lambda - \frac{\hat{u}(s)^2}{3} \right) (1 - \hat{m}_{K_1}^2 - \hat{s}) + 4\hat{m}_\ell^2\lambda \right) \right]
$$

$$
- 2\hat{m}_\ell^2 \lambda \left[ \text{Re} \left( \mathcal{F}^{K_1} \mathcal{H}^{K_1*} \right) - \text{Re} \left( \mathcal{G}^{K_1} \mathcal{H}^{K_1*} \right) (1 - \hat{m}_{K_1}^2) \right] + \frac{\hat{m}_\ell^2}{m_{K_1}^2} \hat{s} \lambda |\mathcal{H}^{K_1}|^2 \right\}.
$$

The differential decay rates $d\mathcal{B}(B^- \to K_1^-\mu^+\mu^-)/ds \equiv \tau_{B^-} \times d\Gamma(B \to K_1^-\mu^+\mu^-)/ds$ are plotted in Fig. 1. To illustrate the dependence on $\theta_{K_1}$, we plot the distributions for the differential decay rates with $\theta_{K_1} = -34^\circ$, $-45^\circ$ and $-57^\circ$, respectively. The effects of charmonium resonances become large for the large region with $s \gtrsim 5\text{ GeV}^2$. We find that in the low $s$ region, where $s \approx 2\text{ GeV}^2$, the differential decay rate for $B \to K_1(1400)\mu^+\mu^-$ with $\theta_{K_1} = -57^\circ$ is enhanced by about 80% compared with that with $\theta_{K_1} = -34^\circ$, whereas the rates for $B \to K_1(1270)\mu^+\mu^-$ is not so sensitive to variation of $\theta_{K_1}$. One should note that the distribution in the low $s$ region is dominated by the $1/s$ term arising from $B \to K_1\gamma$; for instance, for the $B \to K_1(1270)\mu^+\mu^-$ decay, it results in the peak at $s \sim 4\hat{m}_\ell^2$ (or exactly at $s = 0$) and contributes about $-30\%$ at around $s = 2\text{ GeV}^2$ for $-57^\circ < \theta_{K_1} < -34^\circ$.

Furthermore, the value of $\theta_{K_1}$ can be well determined from the following ratio of the distributions,

$$
R_{d\Gamma/ds,\mu} \equiv \frac{d\Gamma(B^- \to K_1^-(1400)\mu^+\mu^-)/ds}{d\Gamma(B^- \to K_1^- (1270)\mu^+\mu^-)/ds}.
$$

(38)
FIG. 1: The dilepton invariant mass distributions for differential decay rates $d\mathcal{B}(B^- \rightarrow K_1^- \mu^+ \mu^-)/ds$ in the SM. The central values of inputs are used. The solid, dotted and dashed curves correspond to $\theta_{K_1} = -34^\circ$, $-45^\circ$, $-57^\circ$, respectively. The thick (blue) [thin (red)] curves correspond to values with [without] resonant corrections.

In Fig. 2 we plot the $R_{d\Gamma/ds,\mu}$ as a function of $s$, which is highly insensitive to the resonance contributions and form factors. When the magnitude of $\theta_{K_1}$ is increased, this ratio peaks at about $s = 1.5 \text{ GeV}^2$ (for $\theta_{K_1} \gtrsim 40^\circ$).

B. Branching fractions

In Table VI we summarize the predictions for branching fractions corresponding to $\theta_{K_1} = -(34 \pm 13)^\circ$. The branching fractions for $B \rightarrow K_1\ell^+\ell^-$ and $B \rightarrow K_1\mu^+\mu^-$ are close to $B \rightarrow K^*(892)\ell^+\ell^-$, $B \rightarrow K^*(892)\mu^+\mu^-$ given in [15]. On the other hand, the branching fractions for $B \rightarrow K_1\tau^+\tau^-$ decays are very small since the allowed phase space is quite narrow. In Fig. 3 we plot the non-resonant branching fractions $\mathcal{B}_{nr}(B \rightarrow K_1(1270)\ell^+\ell^-)$ as functions of $\theta_{K_1}$. For the range of $\theta_{K_1} = -(34 \pm 13)^\circ$, we obtain $\mathcal{B}_{nr}(B \rightarrow K_1(1270)\ell^+\ell^-) \gg$
FIG. 2: The ratio of the decay distributions, $R_{d\Gamma/ds,\mu}$ (see the text), as a function of the dimuon invariant mass $s$. The legends are the same as in Fig. 1.

![Graph showing the ratio $R_{d\Gamma/ds,\mu}$ as a function of $s$ in GeV$^2$.]

FIG. 3: Non-resonant branching fractions $B_{nr}(B^- \rightarrow K_1^- \ell^+ \ell^-)$ as functions of $\theta_{K_1}$. (a) The thick solid, thick dashed, thin solid and thin dashed curves correspond to the decays $B \rightarrow K_1(1270)\ell^+ \ell^-$, $K_1(1400)\ell^+ \ell^-$, $K_1(1270)\mu^+ \mu^-$ and $K_1(1400)\mu^+ \mu^-$, respectively. (b) The solid and dashed curves correspond to $B \rightarrow K_1(1270)\tau^+ \tau^-$, $B \rightarrow K_1(1400)\tau^+ \tau^-$, respectively. The vertical lines indicate the allowed range of $\theta_{K_1}$ given in Eq. (2) [26].

![Graph showing $B_{nr}(B^- \rightarrow K_1^- \ell^+ \ell^-)$ as a function of $\theta_{K_1}$ in degrees.]

$B_{nr}(B \rightarrow K_1(1400)\ell^+ \ell^-)$. It should be helpful to define the ratio,

$$R_{\ell,nr} \equiv \frac{B_{nr}(B \rightarrow K_1(1400)\ell^+ \ell^-)}{B_{nr}(B \rightarrow K_1(1270)\ell^+ \ell^-)}.$$  \hspace{1cm} (39)

We show $R_{\ell,nr}$ as functions of the $\theta_{K_1}$ in Fig. 4. These ratios sensitively depend on $\theta_{K_1}$, and are smaller than 0.15 for $-47^\circ \leq \theta_{K_1} \leq -21^\circ$. We predict

$$R_{e,nr} = 0.04^{+0.01+0.11}_{-0.01-0.02}, \quad R_{\mu,nr} = 0.03^{+0.01+0.09}_{-0.01-0.01}, \quad R_{\tau,nr} = 0.02^{+0.01+0.07}_{-0.00-0.02};$$  \hspace{1cm} (40)
TABLE VI: Predictions for the non-resonant branching fractions $\mathcal{B}_{nr}(B \to K_1 \ell^+ \ell^-)$. The first and second errors come from the uncertainty of the form factors and of the $\theta_{K_1}$ within the allowed region [26], respectively.

| Mode                | $\mathcal{B}_{nr} \times 10^6$ | Mode                | $\mathcal{B}_{nr} \times 10^6$ |
|---------------------|--------------------------------|---------------------|--------------------------------|
| $B^- \to K_1^-(1270)e^+e^-$ | $2.7^{+1.5+0.0}_{-1.2-0.3}$ | $B^- \to K_1^+(1270)e^+e^-$ | $2.5^{+1.4+0.0}_{-1.1-0.3}$ |
| $B^- \to K_1^-(1270)\mu^+\mu^-$ | $2.3^{+1.3+0.0}_{-1.0-0.2}$ | $B^- \to K_1^+(1270)\mu^+\mu^-$ | $2.1^{+1.2+0.0}_{-0.9-0.2}$ |
| $B^- \to K_1^-(1270)\tau^+\tau^-$ | $0.08^{+0.04+0.00}_{-0.03-0.01}$ | $B^- \to K_1^+(1270)\tau^+\tau^-$ | $0.08^{+0.04+0.00}_{-0.03-0.01}$ |
| $B^- \to K_1^-(1400)e^+e^-$ | $0.10^{+0.03+0.25}_{-0.03-0.05}$ | $B^- \to K_1^+(1400)e^+e^-$ | $0.09^{+0.03+0.23}_{-0.03-0.04}$ |
| $B^- \to K_1^-(1400)\mu^+\mu^-$ | $0.06^{+0.02+0.18}_{-0.01-0.02}$ | $B^- \to K_1^+(1400)\mu^+\mu^-$ | $0.06^{+0.02+0.18}_{-0.01-0.02}$ |
| $B^- \to K_1^-(1400)\tau^+\tau^-$ | $0.001^{+0.000+0.005}_{-0.000-0.001}$ | $B^- \to K_1^+(1400)\tau^+\tau^-$ | $0.001^{+0.000+0.005}_{-0.000-0.001}$ |

FIG. 4: $R_{\ell,nr} \equiv \mathcal{B}_{nr}(B \to K_1(1400)\ell^+\ell^-)/\mathcal{B}_{nr}(B \to K_1(1270)\ell^+\ell^-)$ as functions of $\theta_{K_1}$. The solid, dashed and dot-dashed curves correspond to $R_{e,nr}$, $R_{\mu,nr}$ and $R_{\tau,nr}$, respectively. The allowed range of $\theta_{K_1}$ given in Eq. [2] [26] is also shown.

where the first and second errors correspond to the uncertainties of the form factors and $\theta_{K_1}$, respectively. In Fig. 6 we will further show that the ratio $R_{\mu,nr}$ is highly insensitive to the NP corrections.
C. Forward-backward asymmetry

The differential forward-backward asymmetry of the $\overline{B} \to \overline{K}_1 \ell^+ \ell^-$ decay is defined by

$$\frac{dA_{FB}}{d\hat{s}} = \int_{0}^{\hat{s}(\hat{s})} d\hat{u} d^2\hat{\Gamma} / d\hat{u} d\hat{s},$$

which can be written in terms of quantities in Eqs. (29)-(36) as

$$\frac{dA_{FB}}{d\hat{s}} = -\frac{G_F^2 \alpha_{em} \mu_B^5}{2 \pi^5} |V_{ts} V_{tb}| \hat{s} \hat{u}(\hat{s})^2 \left\{ \text{Re} \left( B^{K_1} \mathcal{E}_{K_1} \right) + \text{Re} \left( A^{K_1} \mathcal{F}_{K_1} \right) \right\},$$

and, after including the hard spectator correction [16], are given by

$$\frac{dA_{FB}}{d\hat{s}} = -\frac{G_F^2 \alpha_{em} \mu_B^5}{2 \pi^5} |V_{ts} V_{tb}| \hat{s} \hat{u}(\hat{s})^2 \times c_{10} \left[ \text{Re} (c_9 \hat{s}) A^{K_1} V_1^{K_1} T_1^{K_1} (1 - \hat{m}_{K_1}) + \hat{m}_b \phi \Delta_{HS} \right],$$

where $\Delta_{HS}$ is the hard spectator correction given by

$$\Delta_{HS} = \left\{ (1 + \hat{m}_{K_1}) V_1^{K_1} + (1 - \hat{m}_{K_1})(1 - \hat{s}) A^{K_1} \right\} \times \alpha_s(\mu_B) C_F \pi^2 \int_{0}^{1} d\hat{u} \Phi_{K_1}(u) T_{1,1}^{(n)}(u).$$

Here $f_{K_1}^{\perp}$ and $\Phi_{K_1}^{\perp}(u)$ are the transverse decay constant and the twist-2 tensor light-cone distribution amplitude of the $K_1$, respectively. $f_{K_1(1270)}^{\perp}$, $f_{K_1(1400)}^{\perp}$ and $\Phi_{K_1(1270)}^{\perp}$, $\Phi_{K_1(1400)}^{\perp}$ are related with $f_{K_1A}^{\perp}$, $f_{K_1B}^{\perp}$ and $\Phi_{K_1A}^{\perp}$, $\Phi_{K_1B}^{\perp}$ by [35]

$$\begin{pmatrix} f_{K_1(1270)}^{\perp} \\ f_{K_1(1400)}^{\perp} \end{pmatrix} = M \cdot \begin{pmatrix} f_{K_1A(1270)}^{\perp} \\ f_{K_1B(1270)}^{\perp} \end{pmatrix},$$

$$\begin{pmatrix} f_{K_1(1270)}^{\perp} \\ f_{K_1(1400)}^{\perp} \end{pmatrix} = M \cdot \begin{pmatrix} f_{K_1A(1400)}^{\perp} \\ f_{K_1B(1400)}^{\perp} \end{pmatrix},$$

where $\Phi_{K_1A}^{\perp}$ and $\Phi_{K_1B}^{\perp}$ are expanded as

$$\Phi_{K_1A}^{\perp}(u) = 6u\bar{u} \left[ a_0^{K_{1A,\perp}} + 3 a_1^{K_{1B,\perp}} \xi + a_2^{K_{1B,\perp}} \frac{3}{2} (5\xi^2 - 1) \right],$$

$$\Phi_{K_1B}^{\perp}(u) = 6u\bar{u} \left[ 1 + 3 a_1^{K_{1B,\perp}} \xi + a_2^{K_{1B,\perp}} \frac{3}{2} (5\xi^2 - 1) \right],$$

with $a_0^{K_{1B,\perp}} \equiv 1$, $\bar{u} \equiv 1 - u$ and $\xi \equiv u - \bar{u}$. The values of $f_{K_1A}^{\perp}$, $f_{K_1B}^{\perp}$ and the Gegenbauer moments, $a_i^{(K_1,\perp)}$, are tabulated in Table VI.
In the following, to compare the theoretical predictions with the data, we use the normalized differential forward-backward asymmetry as

$$dA_{FB} / ds = dA_{FB} / ds / \Gamma_s.$$  (49)

In Fig. 5, the normalized differential forward-backward asymmetries $dA_{FB}(B^- \rightarrow K_1^- \mu^+ \mu^-)/ds$ versus $s$ are plotted. For $B \rightarrow K_1(1270)\mu^+ \mu^-$ decays, the dependence of $dA_{FB}/ds$ on $\theta_{K_1}$ is negligibly small. For $\theta_{K_1} \lesssim -45^\circ$, $dA_{FB}(B \rightarrow K_1(1400)\mu^+ \mu^-)/ds$ almost vanishes in the region below the $J/\psi$ resonance. We define $s_{K_1}^{0}$ to be the position of zero of the FBA. $s_{K_1}^{0}$ satisfies

$$\text{Re}(c_{eff}^{HS}(\hat{s}_{K_1}^{0})) = m_b \left\{ T_2^{K_1}(\hat{s}_{K_1}^{0}) (1 - \hat{m}_{K_1}) + T_1^{K_1}(\hat{s}_{K_1}^{0}) (1 + \hat{m}_{K_1}) \right\},$$  (50)

which is negative. Here

$$c_{eff}^{HS}(\hat{s}) = c_7^{eff} + \frac{\Delta_{HS}(\hat{s})}{A^{K_1}(\hat{s}) T_2^{K_1}(\hat{s}) (1 - \hat{m}_{K_1}) + V_1^{K_1}(\hat{s}) T_1^{K_1}(\hat{s}) (1 + \hat{m}_{K_1})}. \quad (51)$$

The position of zero appears below the $J/\psi$-resonance region and depends weakly on $\theta_{K_1}$, especially for $B \rightarrow K_1(1270)\ell^+ \ell^-$ as shown in Fig. 5. We obtain the positions of the zeros of forward-backward asymmetries to be

$$s_{K_1}^{0(1270)} = 2.27^{+0.04+0.01}_{-0.07-0.01} \text{ GeV}^2 \quad \text{and} \quad s_{K_1}^{0(1400)} = 2.80^{+0.23+0.74}_{-0.29-0.07} \text{ GeV}^2,$$  (52)

where the first and second errors correspond to the uncertainties of the form factors and $\theta_{K_1} (= -(34 \pm 13)^\circ)$, respectively. In the following section, we will show that, as the $B \rightarrow K^*(892)\ell^+ \ell^-$ decay, for the $B \rightarrow K_1(1270)\ell^+ \ell^-$ decay the position of the zero of the FBA can be a good observable for searching for new-physics effects.

V. NP EFFECTS

In this section, we study the NP corrections to the $B^- \rightarrow K_1^- \mu^+ \mu^-$ decays in the model-independent way. As in Ref. [15], we parametrize the NP contributions to the Wilson coefficients as

$$c_i \equiv c_i^{SM} + c_i^{NP} = R_i c_i^{SM} \quad \text{for} \quad c_i = c_7^{eff}, c_9, c_{10},$$  (53)
FIG. 5: Normalized differential forward-backward asymmetries:
(a) $dA_{FB}(B^- \rightarrow K_1^-(1270)\mu^+\mu^-)/ds$ and (b) $dA_{FB}(B^- \rightarrow K_1^-(1400)\mu^+\mu^-)/ds$. The legends are the same as Fig. 1.

at the scale $\overline{m}_t(\overline{m}_b)$. For simplicity, we assume all $R_i$ are real. The model-independent analysis for $B \rightarrow X_s\gamma$ and $B \rightarrow X_s\ell^+\ell^-$ [17] gives the following constraints,

$$0.8 \lesssim |R_7| \lesssim 1.2, \quad 1 \lesssim R_9^2 + R_{10}^2 \lesssim 4.$$  

(54)

The possibility of flipped sign of $c_{7\text{eff}}$ due to the NP contribution in the minimum supersymmetric standard model (MSSM) with the minimal flavor violation (MFV) ansatz and with large $\tan\beta$ has been studied in Ref. [18]. The twofold constraint was given by

$$-0.02 \leq c_{7\text{NP}} \leq 0.12 \quad \text{or} \quad 0.59 \leq c_{7\text{NP}} \leq 1.24,$$

(55)

at the weak scale. Further constraints on $c_{7\text{NP}}$ have been obtained with

$$c_{7\text{NP}} = -0.039 \pm 0.043 \cup 0.931 \pm 0.016 \quad (68\% \text{ CL})$$

(56)

$$= [-0.104, 0.026] \cup [0.874, 0.988] \quad (95\% \text{ CL})$$

(57)

in Ref. [37] and

$$c_{7\text{NP}} = 0.02 \pm 0.047 \cup 0.958 \pm 0.002 \quad (68\% \text{ CL})$$

(58)

$$= [-0.039, 0.08] \cup [0.859, 1.031] \quad (95\% \text{ CL})$$

(59)

in Ref. [38]. The sign of $\text{Re}(c_{7\text{eff}})$ can also be flipped in supersymmetric models with non-minimal flavor violation via gluino-down-squark loops. Furthermore, in general flavor-violating supersymmetric models the sign of $c_9$ and $c_{10}$ can be flipped. Therefore, in the
FIG. 6: $R_{\mu, nr} = \mathcal{B}_{nr}(B \to K_1(1400)\mu^+\mu^-)/\mathcal{B}_{nr}(B \to K_1(1270)\mu^+\mu^-)$ as a function of $\theta_{K_1}$. Variations of NP with $(R_7, R_9, R_{10}) = (r, 1, 1), (1, r, 1)$, and $(1, 1, r)$ are respectively included, where $r = 1.0$ (solid), 1.2 (dotted), 0.8 (dot-dashed) and $-1.0$ (dashed). The vertical lines indicate the allowed range of $\theta_{K_1}$ given in Eq. (2) [26].

In Fig. 6, the ratio of the non-resonant branching fractions $R_{\mu, nr} \equiv \mathcal{B}_{nr}(B \to K_1(1400)\mu^+\mu^-)/\mathcal{B}_{nr}(B \to K_1(1270)\mu^+\mu^-)$, including the NP corrections, as a function of the value of $\theta_{K_1}$ is depicted. We show that $R_{\mu, nr}$ is highly insensitive to the NP effect and thus is suitable for determining the value of $\theta_{K_1}$. In Fig. 7, we plot $R_{d\Gamma/ds, \mu}$, the ratio of the differential decay rates, as a function of the dimuon invariant mass, $s$, where the NP effects are considered. We find that $R_{d\Gamma/ds, \mu}$ is insensitive to variation of $R_{10}$, whereas its value is increased (decreased) by about 100% (40%) at about $s = 1.5$ GeV$^2$ corresponding to $\theta_{K_1} = -34^0 (-57^0)$ when $R_7$ or $R_9$ equals to $-1$.

Taking into account the possible NP corrections, we plot $dA_{FB}(B^- \to K_1^-(1270)\mu^+\mu^-)/ds$ as a function of $s$ in Fig. 8. We do not consider the $B^- \to K_1^-(1400)\mu^+\mu^-$ decay, since its branching fraction is relatively small. As shown in Fig. 5 (see also Fig. 8), the differential forward-backward asymmetry for $B \to K_1(1270)\mu^+\mu^-$ and its $s_{0}^{K_1(1270)}$ (if existing) are very insensitive to variation of $\theta_{K_1}$. For the cases with $c_7^{eff}, c_9$ and $c_{10}$ of SM-like sign, the change of the FBA zero owing to variation of NP parameters could be manifest as compared to the
FIG. 7: $R_{d\Gamma/ds,\mu}$, the ratio of the differential decay rates, as a function of the dimuon invariant mass, $s$. Variations of $R_7$, $R_9$ and $R_{10}$ are depicted in (a), (b) and (c) respectively, where the remaining $R_i$ are set to their SM values. The thick (blue) and thin (red) curves correspond to $\theta_{K_1} = -34^\circ$ and $-57^\circ$, respectively. The solid, dotted, dot-dashed and dashed curves correspond to $R_i = 1.0, 1.2, 0.8$ and $-1.0$, respectively.

hadronic uncertainties. As is well known in the case of $B \to K^*(892)\ell^+\ell^-$, for the flipped sign of $c_7^{\text{eff}}$, $c_9$ or $c_{10}$ the characteristic features of the FBA change dramatically. Because the asymmetry zero exists only for $\text{Re}(c_9^{\text{eff}})/c_7^{\text{eff}} < 0$ (see Eq. (50)), therefore there is no asymmetry zero for $(R_7, R_9) = (\pm 1, \mp 1)$ in the spectrum. Flipping the sign of $c_{10}$ would change the sign of the FBA. From the above discussions we can conclude that the position of the FBA zero for the $B \to K_1(1270)\ell^+\ell^-$ decay is a suitable quantity to constrain the NP parameters. Recent measurements for $B \to K^*(892)\ell^+\ell^-$ decays [8, 9] seem to favor (i) the flipped sign of $c_7^{\text{eff}}$ which is denoted by the dashed curves in Fig. (8a), or (ii) the simultaneous flip of the sign of $c_9$ and $c_{10}$ which are denoted by the double-dot dashed curves in Fig. (8c). However, they disfavor the flipped sign($c_9c_{10}$) models. See also the discussion in Ref. [20].
FIG. 8: Normalized differential forward-backward asymmetry $dA_{FB}(B^- \to K_1^-(1270)\mu^+\mu^-)/ds$ as a function of the dimuon invariant mass $s$. The thick (blue) and thin (red) curves correspond to the asymmetries with $\theta_{K_1} = -34^\circ$ and $-57^\circ$, respectively. In (a), where $R_9 = R_{10} = 1.0$ (the SM result), the solid curves are for $R_7 = 1.0$, the dotted for $R_7 = 1.2$, the dot-dashed for $R_7 = 0.8$ and the dashed for $R_7 = -1.0$. In (b), where $R_7 = R_{10} = 1.0$ (the SM result), and the solid curves are for $R_9 = 1.0$, the dotted for $R_9 = 1.2$, the dot-dashed for $R_9 = 0.8$ and the dashed for $R_9 = -1.0$. In (c), where $R_7 = R_9 = 1.0$ and the solid curves are for $R_{10} = 1.0$ (the SM result), the dotted for $R_{10} = 1.2$, the dot-dashed for $R_{10} = 0.8$ and the dashed for $R_{10} = -1.0$. The $dA_{FB}/ds$ with $(R_7, R_9, R_{10}) = (1.0, -1.0, -1.0)$, and $(-1.0, 1.0, -1.0)$ are denoted by the double-dot dashed and long-short dashed curves, respectively, in (c).

VI. SUMMARY

We have studied the rare decays $B \to K_1\ell^+\ell^-$ with $K_1 \equiv K_1(1270)$, $K_1(1400)$ and $\ell \equiv e, \mu, \tau$. The strange axial-vector mesons, $K_1(1270)$ and $K_1(1400)$, are the mixtures of the $K_{1A}$ and $K_{1B}$, which are the $1^3P_1$ and $1^1P_1$ states, respectively. Although the branching ratios depend on the magnitudes of $B \to K_1$ form factors, the $K_1(1270)$-
$K_1(1400)$ mixing angle, $\theta_{K_1}$, can be well determined from the measurement of the ratio $R_\ell \equiv \mathcal{B}(B \to K_1(1400)\ell^+\ell^-)/\mathcal{B}(B \to K_1(1270)\ell^+\ell^-)$, which depends very weakly on new-physics corrections. We have calculated differential forward-backward asymmetries of $B \to K_1\ell^+\ell^-$ decays. For $B \to K_1(1270)\mu^+\mu^-$, the asymmetry zero, which depends very weakly on $\theta_{K_1}$, can be dramatically changed due to variation of new-physics parameters.

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[1] E. Barberio et al., arXiv:0808.1297 [hep-ex].
[2] B. Aubert et al. [BABAR Collaboration], Phys. Rev. D 70, 112006 (2004) [arXiv:hep-ex/0407003].
[3] M. Nakao et al. [Belle Collaboration], Phys. Rev. D 69, 112001 (2004) [arXiv:hep-ex/0402042].
[4] T. E. Coan et al. [CLEO Collaboration], Phys. Rev. Lett. 84, 5283 (2000) [arXiv:hep-ex/9912057].
[5] H. Yang et al., Phys. Rev. Lett. 94, 111802 (2005) [arXiv:hep-ex/0412039].
[6] A. Ishikawa et al. [Belle Collaboration], Phys. Rev. Lett. 91, 261601 (2003) [arXiv:hep-ex/0308044].
[7] B. Aubert et al. [BABAR Collaboration], Phys. Rev. D 73, 092001 (2006) [arXiv:hep-ex/0604007].
[8] A. Ishikawa et al. [Belle Collaboration], Phys. Rev. Lett. 96, 251801 (2006) [arXiv:hep-ex/0603018].
[9] B. Aubert et al. [BABAR Collaboration], arXiv:0804.4412 [hep-ex].
[10] B. Aubert et al. [BABAR Collaboration], arXiv:0807.4119 [hep-ex].
[11] G. Eigen, arXiv:0807.4076 [hep-ex].
[12] M. A. Paracha, I. Ahmed and M. J. Aslam, Eur. Phys. J. C 52, 967 (2007) arXiv:0707.0733 [hep-ph].
[13] I. Ahmed, M. A. Paracha and M. J. Aslam, Eur. Phys. J. C 54, 591 (2008) [arXiv:0802.0740 [hep-ph]].

[14] A. Saddique, M. J. Aslam and C. D. Lu, arXiv:0803.0192 [hep-ph].

[15] A. Ali, P. Ball, L. T. Handoko and G. Hiller, Phys. Rev. D 61, 074024 (2000) [arXiv:hep-ph/9910221].

[16] M. Beneke, T. Feldmann and D. Seidel, Nucl. Phys. B 612, 25 (2001) [arXiv:hep-ph/0106067].

[17] A. Ali, E. Lunghi, C. Greub and G. Hiller, Phys. Rev. D 66, 034002 (2002) [arXiv:hep-ph/0112300].

[18] T. Feldmann and J. Matias, JHEP 0301, 074 (2003) [arXiv:hep-ph/0212158].

[19] F. Kruger and J. Matias, Phys. Rev. D 71, 094009 (2005) [arXiv:hep-ph/0502060].

[20] C. Bobeth, G. Hiller and G. Piranishvili, JHEP 0807, 106 (2008) [arXiv:0805.2525 [hep-ph]].

[21] U. Egede, T. Hurth, J. Matias, M. Ramon and W. Reece, arXiv:0807.2589 [hep-ph].

[22] C. H. Chen, C. Q. Geng and L. Li, arXiv:0808.0127 [hep-ph].

[23] M. Suzuki, Phys. Rev. D 47, 1252 (1993).

[24] L. Burakovsky and J. T. Goldman, Phys. Rev. D 57, 2879 (1998) [arXiv:hep-ph/9703271].

[25] H. Y. Cheng, Phys. Rev. D 67, 094007 (2003) [arXiv:hep-ph/0301198].

[26] H. Hatanaka and K. C. Yang, Phys. Rev. D 77, 094023 (2008) [arXiv:0804.3198 [hep-ph]].

[27] A. J. Buras, M. Misiak, M. Munz and S. Pokorski, Nucl. Phys. B 424, 374 (1994) [arXiv:hep-ph/9311345].

[28] A. J. Buras and M. Munz, Phys. Rev. D 52, 186 (1995) [arXiv:hep-ph/9501281].

[29] C. S. Lim, T. Morozumi and A. I. Sanda, Phys. Lett. B 218, 343 (1989).

[30] A. Ali, T. Mannel and T. Morozumi, Phys. Lett. B 273, 505 (1991).

[31] F. Kruger and L. M. Sehgal, Phys. Lett. B 380, 199 (1996) [arXiv:hep-ph/9603237].

[32] W. M. Yao et al. [Particle Data Group], J. Phys. G 33, 1 (2006).

[33] K. C. Yang, Phys. Rev. D 78, 034018 (2008) [arXiv:0807.1171 [hep-ph]].

[34] K. C. Yang, in preparation.

[35] K. C. Yang, Nucl. Phys. B 776, 187 (2007) [arXiv:0705.0692 [hep-ph]].

[36] CKM fitter group (URL: http://ckmfitter.in2p3.fr).

[37] U. Haisch and A. Weiler, Phys. Rev. D 76, 074027 (2007) [arXiv:0706.2054 [hep-ph]].

[38] C. Bobeth, M. Bona, A. J. Buras, T. Ewerth, M. Pierini, L. Silvestrini and A. Weiler, Nucl.
Phys. B 726, 252 (2005) [arXiv:hep-ph/0505110].