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# $B \to K_2 \ell^+ \ell^-$ decay beyond the Standard Model

T. M. Aliev *†, M. Savcı ‡

Physics Department, Middle East Technical University, 06531 Ankara, Turkey

**Abstract**

The exclusive $B \to K_2 \ell^+ \ell^-$ decay is studied using the most general, model independent four–fermion interaction. The sensitivity of the ratio of the decay widths when $K_2$ meson is longitudinally and transversally polarized, the forward–backward asymmetry and longitudinal polarization of the final lepton on the new Wilson coefficients is studied. It is found that these quantities are very useful for establishing new physics beyond the Standard Model.

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*e-mail: taliev@metu.edu.tr

†permanent address:Institute of Physics,Baku,Azerbaijan

‡e-mail: savci@metu.edu.tr

# 1 Introduction

Flavor changing neutral current (FCNC) processes provide promising direction for testing the gauge structure of the Standard Model (SM) at loop level, as these decays are forbidden at tree level. Moreover, FCNC decays are sensitive to the new physics beyond the SM. Among all the FCNC processes the rare $B$–meson decays occupy an important place since they contain rich phenomena relevant to the SM and new physics beyond it. The rare decays due to $b \to s$ transition receives special attention since SM predicts "large" branching ratio for them.

The radiative $B \to K^*\gamma$ [1–3], $B \to K_1(1270, 1430)\gamma$ [4] and semileptonic $B \to K^*(892)\ell^+\ell^-$ [5, 6] have been measured in experiments. [4]. The isospin and forward–backward asymmetry in $B \to K^*(892)\ell^+\ell^-$ is also measured by BaBar Collaboration [7, 8]. The semileptonic decays $B \to K(K^*)\ell^+\ell^-$ theoretically are studied intensively in literature [11]. The radiative $B \to K_2(1430)\gamma$ decay has also been observed at BaBar [9, 10].

The exclusive $B \to K_2(1430)\ell^+\ell^-$ decay is studied within the SM in [12, 13] and it is obtained that the branching ratio of this channel is comparable with that of the $B \to K^*(892)\ell^+\ell^-$ decay. Therefore, investigation of the $B \to K_2(1430)\ell^+\ell^-$ decay can provide an independent test of the SM. The main similarity of the $B \to K^*(892)\ell^+\ell^-$ and $B \to K_2(1430)\ell^+\ell^-$ decays in the SM is that both decays are described by the $b \to s$ transition, hence by the same effective Hamiltonian.

In SM the $b \to s\ell^+\ell^-$ transition is described by the Wilson coefficients $C_7$, $C_9$ and $C_{10}$ of the operators $\mathcal{O}_7$, $\mathcal{O}_9$ and $\mathcal{O}_{10}$ at $\mu = m_b$ scale. Explicit expressions of these Wilson coefficients and relevant operators can be found in [14]. The new physics in FCNC transitions can appear in two different ways: a) By new contributions to the Wilson coefficients, i.e., by modification of the Wilson coefficients, b) appearance of new operators which are absent in the SM.

It should be noted that, the most general model independent analysis of the $b \to s\ell^+\ell^-$ transition is carried out in [15] in terms of ten new types of local four–Fermi interaction. Furthermore, extensive study on various observables of the $b \to s\ell^+\ell^-$ transition is performed in the same framework in [16, 17]. Extension to the exclusive $B \to K(K^*)\ell^+\ell^-$ channels are performed in [18–20].

One important quantity in checking the predictions of the SM and establishing new physics is the lepton polarization effects, which are first pointed out in [21] and subsequently considered in many works (see for example references in [22]). The main goal of the present work is to study the branching ratio, forward–backward asymmetry and the final lepton polarization effects for the rare $B \to K_2(1430)\ell^+\ell^-$decay in a model independent manner. It should be noted here that, $B \to K_2(\to K\pi)\ell^+\ell^-$ decay is studied in the SM and in the two new physics scenarios, namely vector–like quark model and the familiar non–universal $Z'$ model in [23]. The work is arranged as follows. In section 2, we firstly present the most general form of the local four–Fermi interactions and then using this form we calculate the helicity amplitudes for the $B \to K_2(1430)\ell^+\ell^-$ decay. In this section we also present the expressions of the branching ratio, forward–backward asymmetry and lepton polarizations in terms of helicity amplitudes. Section 3 contains our numerical analysis and conclusions.

## 2 Helicity amplitudes for the $B \to K_2 \ell^+ \ell^-$ decay

The $B \to K_2 \ell^+ \ell^-$ decay is described at quark level by the $b \to s\ell^+\ell^-$ transition. After integrating over heavy degrees of freedom, the effective Hamiltonian in the SM for the $b \to s\ell^+\ell^-$ transition can be expressed as,

$$\mathcal{H}_{eff} = -\frac{4G_F}{\sqrt{2}} V_{tb} V_{ts}^* \sum_{i=1}^{10} C_i(\mu) \mathcal{O}_i(\mu) \ , \tag{1}$$

where $\mathcal{O}_i(\mu)$ are the four–quark operators and $C_i(\mu)$ are the corresponding Wilson coefficients at $\mu$ scale. Explicit expressions of the Wilson coefficients at next–to–next leading logarithm (NNLL) are calculated in many works (for example see [24] and the references therein). The operators responsible for for the $B \to K_2 \ell^+ \ell^-$ decay are $\mathcal{O}_7$, $\mathcal{O}_9$ and $\mathcal{O}_{10}$ are given as,

$$\begin{aligned}
\mathcal{O}_7 &= \frac{e^2}{16\pi^2} m_b (\bar{s}_R \sigma_{\mu\nu} b_R)\, F^{\mu\nu} \ , \\
\mathcal{O}_9 &= \frac{e^2}{16\pi^2} (\bar{s}_L \gamma_\mu b_L)\, \bar{\ell}\gamma_\mu \ell \ , \\
\mathcal{O}_{10} &= \frac{e^2}{16\pi^2} (\bar{s}_L \gamma_\mu b_L)\, \bar{\ell}\gamma_\mu \gamma_5 \ell \ ,
\end{aligned} \tag{2}$$

where $q_{L(R)} = \dfrac{1 \mp \gamma_5}{2} q$. Using this effective Hamiltonian, the matrix element for the $b \to s\ell^+\ell^-$ transition can be written as,

$$\mathcal{M} = \frac{G_F}{\sqrt{2}\pi} V_{tb} V_{ts}^* \left\{ C_9^{eff} \bar{s}_L \gamma_\mu b_L \, \bar{\ell}\gamma_\mu \ell + C_{10} \bar{s}_L \gamma_\mu b_L \, \bar{\ell}\gamma_\mu \gamma_5 \ell - \frac{2C_7^{eff}}{q^2} \bar{s}_R i \sigma_{\mu\nu} q^\nu b_R \, \bar{\ell}\gamma_\mu \ell \right\} \ , \tag{3}$$

where $C_9^{eff}$ contains short and long distance contributions and it can be written as,

$$C_9^{eff} = C_9(\mu) + Y_{SD}(z, \hat{s}) + Y_{LD}(z, \hat{s}) \ . \tag{4}$$

Here, $z = m_c/m_b$, $\hat{s} = q^2/m_b^2$, $Y_{SD}$ describes contributions coming from four–quark operators and $Y_{LD}$ corresponds to the long distance effects from four–quark operators near the $\bar{c}c$ resonance. The expressions for $Y_{SD}$ and $Y_{SD}$ can be written as

$$\begin{aligned}
Y_{SD}(z, \hat{s}) = {}& h\,(z, \hat{s})\, C^{(0)}(\mu) - \frac{1}{2} h\,(1, \hat{s})\, [4C_3(\mu) + 4C_4(\mu) + 3C_5(\mu) + C_6(\mu)] \\
& - \frac{1}{2} h\,(0, \hat{s})\, [C_3(\mu) + 3C_4(\mu)] + \frac{2}{9}\, [3C_3(\mu) + C_4(\mu) + 3C_5(\mu) + C_6(\mu)] \ ,
\end{aligned} \tag{5}$$

where $x = 4z^2/\hat{s}$, and,

$$\begin{aligned}
C^{(0)}(\mu) &= 3C_1(\mu) + C_2(\mu) + 3C_3(\mu) + C_4(\mu) + 3C_5(\mu) + C_6(\mu) \ , \\
h\,(z, \hat{s}) &= -\frac{8}{9} \ln z + \frac{8}{27} + \frac{4}{9} x - \frac{2}{9} (2 + x) \sqrt{|1 - x|} \\
& \quad \times \left[ \Theta(1 - x) \left( \ln \left| \frac{1 + \sqrt{1 - x}}{1 - \sqrt{1 - x}} \right| - i\pi \right) + \Theta(x - 1)\, 2 \arctan \frac{1}{\sqrt{x - 1}} \right], \\
h\,(0, \hat{s}) &= -\frac{8}{27} - \frac{8}{9} \ln \frac{m_b}{\mu} - \frac{4}{9} \ln \hat{s} + \frac{4}{9} i\pi \ ,
\end{aligned} \tag{6}$$

and

$$Y_{LD}(\hat{s}) = \frac{3}{\alpha^2} C^{(0)}(\mu) \sum_{V_i=\psi(1s),\cdots,\psi(6s)} \frac{\pi æ_i \Gamma\left(V_i \to \ell^+\ell^-\right) M_{V_i}}{\left(M_{V_i}^2 - \hat{s} m_b^2 - i M_{V_i}\Gamma_{V_i}\right)} , \tag{7}$$

The values of the phenomenological parameters $æ_i$ are fixed from the analysis of $B \to K^*\ell^+\ell^-$ decay and they are taken to be $æ = 1.65$ and $æ = 2.36$ for the resonances $J/\psi$ and $\psi'$, respectively.

The charm loop brings further corrections to the radiative $b \to s\gamma$ transition, which modifies the Wilson coefficient $C_7^{eff}$. The Wilson coefficient $C_7^{eff}$ can be written as [25],

$$C_7^{eff} = C_7(\mu) + C_7'(\mu) ,$$

where

$$\begin{aligned} C_7'(\mu) &= i\alpha_s \left\{ \frac{2}{9}\eta^{14/23} \left[ F(x_t) - 0.1687 \right] - 0.03 C_2(\mu) \right\} , \\ F(x_t) &= \frac{x_t(x_t^2 - 5x_t - 2)}{8(x_t-1)^2} + \frac{3(x_t \ln x_t)^2}{4(x_t-1)^4} , \end{aligned} \tag{8}$$

with $x_t = m_t^2/m_W^2$ and $\eta = \alpha_s(m_W)/\alpha_s(\mu)$.

As has already been noted, our main aim in this work is to analyze the $B \to K_2(1430)\ell^+\ell^-$ decay in a model independent way.The most general, model independent local four–Fermi interaction is given in [15], which might contribute to the considered decay. Explicit form of the local four–Fermi interaction which describes $b \to s\ell^+\ell^-$ transition can be written as,

$$\begin{aligned} \mathcal{M}_{new} = \frac{G_F\alpha}{\sqrt{2}\pi} V_{tb}V_{ts}^* \Bigg\{ & C_{BR}\bar{s}_L i\sigma_{\mu\nu}\frac{q^\nu}{q^2} b_R \bar{\ell}\gamma^\mu\ell + C_{SL}\bar{s}_R i\sigma_{\mu\nu}\frac{q^\nu}{q^2} b_L \bar{\ell}\gamma^\mu\ell + C_{LL}\bar{s}_L\gamma_\mu b_L \,\bar{\ell}\gamma_\mu\ell \\ & + C_{LR}\bar{s}_L\gamma_\mu b_L \bar{\ell}_R\gamma^\mu\ell_R + C_{RL}\bar{s}_R\gamma_\mu b_R \bar{\ell}_L\gamma^\mu\ell_L + C_{RR}\bar{s}_R\gamma_\mu b_R \bar{\ell}_R\gamma^\mu\ell_R \\ & + C_{LRRL}\bar{s}_L b_R \bar{\ell}_R\ell_L + C_{LRLR}\bar{s}_L b_R \bar{\ell}_L\ell_R + C_{RLRL}\bar{s}_R b_L \bar{\ell}_R\ell_L \\ & + C_{RLLR}\bar{s}_R b_L \bar{\ell}_L\ell_R + C_T\bar{s}\sigma_{\mu\nu} b \bar{\ell}\sigma^{\mu\nu}\ell + iC_{TE}\epsilon_{\mu\nu\alpha\beta}\bar{s}\sigma^{\mu\nu} b \bar{\ell}\sigma^{\alpha\beta}\ell \Bigg\} , \end{aligned} \tag{9}$$

For the sake of simplicity we neglect the contribution of the tensor interaction in further discussion.

Using Eqs. (3) and (9), the matrix element of the $b \to s\ell^+\ell^-$ transition, including the SM and new physics contributions, can be written as,

$$\mathcal{M} = \mathcal{M}_{SM} + \mathcal{M}_{new} . \tag{10}$$

The matrix element for the exclusive $B \to K_2(1430)\ell^+\ell^-$ decay can be obtained from Eq. (10) by sandwiching it between initial and final states, i.e.,

$$\left\langle K_2(1430)(p,\varepsilon) \left| \mathcal{M} \right| B(p_B) \right\rangle . \tag{11}$$

Before giving definition of the matrix of quark operators between initial and final meson states, few words about the helicity states of the tensor $K_2(1430)$ meson are in order. The polarizations $\varepsilon_\lambda^{\mu\nu}$ with helicity $\lambda$ of the tensor meson with mass $m$ and four–momentum $(E,0,0,p_z)$ moving along the z–axis can be written in terms of the polarization vectors [26]

$$\begin{aligned}
\varepsilon_{(0)}^{*\mu} &= \frac{1}{m}(p,0,0,E) \ , \\
\varepsilon_{\pm}^{*\mu} &= \frac{1}{\sqrt{2}}(0,\mp 1,i,0) \ ,
\end{aligned} \qquad (12)$$

in the following way,

$$\begin{aligned}
\varepsilon_{\pm 2}^{*\alpha\beta} &= \varepsilon_{\pm}^{\alpha}\varepsilon_{\pm}^{\beta} \ , \\
\varepsilon_{\pm 1}^{*\alpha\beta} &= \frac{1}{\sqrt{2}}\Big[\varepsilon_{\pm}^{\alpha}\varepsilon_{0}^{\beta} + \varepsilon_{0}^{\alpha}\varepsilon_{\pm}^{\beta}\Big] \ , \\
\varepsilon_{0}^{*\alpha\beta} &= \frac{1}{\sqrt{6}}\Big[\varepsilon_{+}^{\alpha}\varepsilon_{-}^{\beta} + \varepsilon_{-}^{\alpha}\varepsilon_{+}^{\beta}\Big] + \sqrt{\frac{2}{3}}\varepsilon_{0}^{\alpha}\varepsilon_{0}^{\beta} \ .
\end{aligned} \qquad (13)$$

It follows from the expression of the amplitude of the $b \to s\ell^+\ell^-$ transition that, in order to obtain the matrix element for the semileptonic $B \to K_2(1430)\ell^+\ell^-$ decay, the following matrix elements are needed to be known,

$$\langle K_2(1430)(p,\varepsilon)\left|\bar{s}\gamma_\mu(1\pm\gamma_5)b\right|B(p_B)\rangle \ , \qquad (14)$$

$$\langle K_2(1430)(p,\varepsilon)\left|\bar{s}\sigma_{\mu\nu}q^\nu(1+\gamma_5)b\right|B(p_B)\rangle \ , \qquad (15)$$

$$\langle K_2(1430)(p,\varepsilon)\left|\bar{s}(1\pm\gamma_5)b\right|B(p_B)\rangle \ . \qquad (16)$$

The matrix element given in Eq. (14) can be parametrized in terms of the form factors as follows,

$$\begin{aligned}
&\langle K_2(p,\varepsilon)\left|\bar{s}\gamma_\mu(1\pm\gamma_5)b\right|B(p_B)\rangle = -\epsilon_{\mu\nu\alpha\beta}\varepsilon^{*\nu}p^\alpha q^\beta \frac{2V(q^2)}{m_B+m_{K_2}} \pm i\varepsilon_\mu^*(m_B+m_{K_2})A_1(q^2) \\
&\mp i(p_B+p)_\mu(\varepsilon^* q)\frac{A_2(q^2)}{m_B+m_{K_2}} \mp iq_\mu\frac{2m}{q^2}(\varepsilon^* q)\left[A_3(q^2)-A_0(q^2)\right] \ ,
\end{aligned} \qquad (17)$$

$$\begin{aligned}
&\langle K_2(p,\varepsilon)\left|\bar{s}\sigma_{\mu\nu}q^\nu(1+\gamma_5)b\right|B(p_B)\rangle = 2\epsilon_{\mu\nu\alpha\beta}\varepsilon^{*\nu}p^\alpha q^\beta T_1(q^2) + i\Big[\varepsilon_\mu^*(m_B^2-m_{K_2}^2) \\
&-(p_B+p)_\mu(\varepsilon^* q)\Big]T_2(q^2) + i(\varepsilon^* q)\left[q_\mu - (p_B+p)_\mu\frac{q^2}{m_B^2-m_{K_2}^2}\right]T_3(q^2) \ ,
\end{aligned} \qquad (18)$$

where,

$$\varepsilon_\lambda^{*\mu} \equiv \varepsilon_\lambda^{*\mu\nu}\frac{q_\nu}{m_B} \ .$$

The matrix element $\langle K_2(p,\varepsilon)\left|\bar{s}(1\pm\gamma_5)b\right|B(p_B)\rangle$ can be obtained from Eq. (17) by multiplying it with $q_\mu$ and then using equation of motion. Neglecting strange quark mass we get,

$$\begin{aligned}
&\langle K_2(p,\varepsilon)\left|\bar{s}(1\pm\gamma_5)b\right|B(p_B)\rangle = \frac{1}{m_B}\Big\{i(\varepsilon^* q)(m_B+m_{K_2})A_1(q^2) \\
&\pm(m_B-m_{K_2})(\varepsilon^* q)A_2(q^2) \pm 2m(\varepsilon^* q)\left[A_3(q^2)-A_0(q^2)\right]\Big\} \ .
\end{aligned} \qquad (19)$$

The following relation holds among the form factors $A_1(q^2)$, $A_2(q^2)$ and $A_3(q^2)$,

$$2mA_3(q^2) = (m_B + m_{K_2})A_1(q^2) - (m_B - m_{K_2})A_2(q^2) \ . \tag{20}$$

Using Eqs.(19) and (20), we get

$$\langle K_2(p,\varepsilon)\left|\bar{s}(1\pm\gamma_5)b\right|B(p_B)\rangle = \frac{1}{m_B}\left[\pm 2m(\varepsilon^* q)A_0(q^2)\right] \ . \tag{21}$$

Using these definitions of the form factors, we get the following expression for the decay amplitude of the $B \to K_2(1430)\ell^+\ell^-$ channel,

$$\begin{aligned}
\mathcal{M} = \frac{G_F\alpha}{4\sqrt{2}\pi}V_{tb}V_{ts}^*\Bigg\{&\bar{\ell}\gamma_\mu(1-\gamma_5)\ell\Big[-2A_1\epsilon_{\mu\nu\lambda\sigma}\varepsilon^{\nu*}p^\lambda q^\sigma - iB_1\varepsilon_\mu^* + iB_2(\varepsilon^* q)(p_B+p)_\mu \\
&+ iB_3(\varepsilon^* q)q_\mu\Big] + \bar{\ell}\gamma_\mu(1+\gamma_5)\ell\Big[-2C_1\epsilon_{\mu\nu\lambda\sigma}\varepsilon^{\nu*}p^\lambda q^\sigma - iD_1\varepsilon_\mu^* + iD_2(\varepsilon^* q)(p_B+p)_\mu \\
&+ iD_3(\varepsilon^* q)q_\mu\Big] + \bar{\ell}(1-\gamma_5)\ell\Big[iB_4(\varepsilon^* q)\Big] + \bar{\ell}(1+\gamma_5)\ell\Big[iB_5(\varepsilon^* q)\Big]\Bigg\} \ .
\end{aligned} \tag{22}$$

Here

$$\begin{aligned}
A_1 &= (C_{LL}^{tot} + C_{RL})\frac{V}{m_B + m_{K_2}} - (C_{BR} + C_{SL})\frac{T_1}{q^2} \ , \\
B_1 &= (C_{LL}^{tot} - C_{RL})(m_B + m_{K_2})A_1 - (C_{BR} - C_{SL})(m_B^2 - m_{K_2}^2)\frac{T_2}{q^2} \ , \\
B_2 &= \frac{C_{LL}^{tot} - C_{RL}}{m_B + m_{K_2}}A_2 - (C_{BR} - C_{SL})\frac{1}{q^2}\left[T_2 + \frac{q^2}{m_B^2 - m_{K_2}^2}T_3\right] \ , \\
B_3 &= 2(C_{LL}^{tot} - C_{RL})m\frac{A_3 - A_0}{q^2} + (C_{BR} - C_{SL})\frac{T_3}{q^2} \ , \\
C_1 &= A_1(C_{LL}^{tot} \to C_{LR}^{tot} \ , \quad C_{RL} \to C_{RR}) \ , \\
D_1 &= B_1(C_{LL}^{tot} \to C_{LR}^{tot} \ , \quad C_{RL} \to C_{RR}) \ , \\
D_2 &= B_2(C_{LL}^{tot} \to C_{LR}^{tot} \ , \quad C_{RL} \to C_{RR}) \ , \\
D_3 &= B_3(C_{LL}^{tot} \to C_{LR}^{tot} \ , \quad C_{RL} \to C_{RR}) \ , \\
B_4 &= -2(C_{LRRL} - C_{RLRL})\frac{m}{m_b}A_0 \ , \\
B_5 &= -2(C_{LRLR} - C_{RLLR})\frac{m}{m_b}A_0 \ ,
\end{aligned} \tag{23}$$

where

$$\begin{aligned}
C_{LL}^{tot} &= C_9^{eff} - C_{10} + C_{LL} \ , \\
C_{LR}^{tot} &= C_9^{eff} + C_{10} + C_{LR} \ , \\
C_{BR} &= -2m_b C_7^{eff} + C'_{BR} \ .
\end{aligned} \tag{24}$$

It should be stressed at this point that, the matrix element of $B \to K_2(1430)\ell^+\ell^-$ decay is formally the same with that of the $B \to V\ell^+\ell^-$ decay ($V$ is $\rho$ or $K$ meson). But it is necessary to keep in mind that form factors in both cases are different, and also, the polarization vector $\varepsilon^*_\alpha$ which has the form,

$$\varepsilon^*_\alpha = \frac{\varepsilon^*_{\alpha\beta} q^\beta}{m_B} ,$$

is different from the polarization vector of the vector mesons.

Having obtained the matrix element $B \to K_2(1430)\ell^+\ell^-$ decay, the next step in our analysis is to calculate the helicity amplitudes for this decay. Using the helicity amplitude formalism presented in [27], we get the following helicity amplitudes for the $B \to K_2(1430)\ell^+\ell^-$ decay,

$$\begin{aligned}
M^{++}_\pm &= \mp i \frac{m_\ell}{m_B m_{K_2}} \left|\vec{p}_{K_2}\right| \sqrt{q^2} \sin\theta \Big[ (B_1 + D_1) \mp 2 \left|\vec{p}_{K_2}\right| \sqrt{q^2}\, (A_1 + C_1) \Big] , \\
M^{+-}_\pm &= i \frac{(\mp 1 + \cos\theta)}{2 m_B m_{K_2}} \left|\vec{p}_{K_2}\right| q^2 \Big\{ \mp (1-v) B_1 \mp (1+v) D_1 \\
&+ 2 \left|\vec{p}_{K_2}\right| \sqrt{q^2} \Big[ (1-v) A_1 + (1+v) C_1 \Big] \Big\} , \\
M^{-+}_\pm &= i \frac{(\pm 1 + \cos\theta)}{2 m_B m_{K_2}} \left|\vec{p}_{K_2}\right| q^2 \Big\{ \mp (1+v) B_1 \mp (1-v) D_1 \\
&+ 2 \left|\vec{p}_{K_2}\right| \sqrt{q^2} \Big[ (1+v) A_1 + (1-v) C_1 \Big] \Big\} , \\
M^{--}_\pm &= \pm i \frac{m_\ell}{m_B m_{K_2}} \left|\vec{p}_{K_2}\right| \sqrt{q^2} \sin\theta \Big[ (B_1 + D_1) \mp 2 \left|\vec{p}_{K_2}\right| \sqrt{q^2}\, (A_1 + C_1) \Big] , \\
M^{++}_0 &= i \frac{\sqrt{2/3}}{m_B m^2_{K_2}} \left|\vec{p}_{K_2}\right| \sqrt{q^2} \Big\{ 2 m_\ell \Big[ E_{K_2} \cos\theta\, (B_1 + D_1) + \left|\vec{p}_{K_2}\right| (B_1 - D_1) \Big] \\
&- 2 m_\ell \left|\vec{p}_{K_2}\right| \sqrt{q^2} \Big[ (B_2 - D_2)\, (E_B + E_{K_2}) + 2 \left|\vec{p}_{K_2}\right| \cos\theta\, (B_2 + D_2) \Big] \\
&- \left|\vec{p}_{K_2}\right| q^2 \Big[ 2 m_\ell\, (B_3 - D_3) + (1+v) B_4 - (1-v) B_5 \Big] \Big\} , \\
M^{+-}_0 &= i \frac{\sqrt{2/3}}{m_B m^2_{K_2}} \left|\vec{p}_{K_2}\right| q^2 \sin\theta \Big\{ - E_{K_2} \Big[ B_1 (1-v) + D_1 (1+v) \Big] \\
&+ 2 \left|\vec{p}_{K_2}\right|^2 \sqrt{q^2} \Big[ B_2 (1-v) + D_2 (1+v) \Big] \Big\} , \\
M^{-+}_0 &= i \frac{\sqrt{2/3}}{m_B m^2_{K_2}} \left|\vec{p}_{K_2}\right| q^2 \sin\theta \Big\{ - E_{K_2} \Big[ B_1 (1+v) + D_1 (1-v) \Big] \\
&+ 2 \left|\vec{p}_{K_2}\right|^2 \sqrt{q^2} \Big[ B_2 (1+v) + D_2 (1-v) \Big] \Big\} , \\
M^{--}_0 &= i \frac{\sqrt{2/3}}{m_B m^2_{K_2}} \left|\vec{p}_{K_2}\right| \sqrt{q^2} \Big\{ - 2 m_\ell \Big[ E_{K_2} \cos\theta\, (B_1 + D_1) + \left|\vec{p}_{K_2}\right| (B_1 - D_1) \Big] \\
&- 2 m_\ell \left|\vec{p}_{K_2}\right| \sqrt{q^2} \Big[ (B_2 - D_2)\, (E_B + E_{K_2}) - 2 \left|\vec{p}_{K_2}\right| \cos\theta\, (B_2 + D_2) \Big]
\end{aligned}$$

$$- \left|\vec{p}_{K_2}\right| q^2 \Big[ 2m_\ell \left(B_3 - D_3\right) + (1-v) B_4 - (1+v) B_5 \Big] \Big\} ~. \tag{25}$$

Here superscripts and subscripts denote the helicities of the leptons and $K_2$ meson, respectively. In Eq. (26) we have,

$$\begin{aligned}
\lambda(m_B^2, q^2, m_{K_2}^2) &= m_B^4 + q^4 + m_{K_2}^4 - 2 m_B^2 q^2 - 2 m_B^2 m_{K_2}^2 - 2 q^2 m_{K_2}^2 ~, \\
q^2 &= (p_B - p_{K_2})^2 = (p_1 + p_2)^2 ~, \\
v &= \sqrt{1 - \frac{4 m_\ell^2}{q^2}} ~, \\
\left|\vec{p}_{K_2}\right| &= \frac{\lambda^{1/2}(m_B^2, q^2, m_{K_2}^2)}{2 m_B} ~,
\end{aligned}$$

and $m_\ell$ is the lepton mass, $\theta$ is the angle between $K_2$ and $\ell^-$ lepton.

It should be noted here that the $\pm 2$ helicity states of the tensor meson give no contribution to the helicity amplitudes. This is due to the fact that in the CM of leptons only time component of $q^2$ is different from zero, and therefore $\varepsilon_\pm^{*\alpha} q_\alpha = 0$.

As a result of some calculation, we obtain the differential decay width in terms of the helicity amplitudes as follows,

$$\begin{aligned}
\frac{d\Gamma}{dq^2\, d\cos\theta} &= \frac{G_F^2 \alpha^2 \left|V_{tb} V_{ts}^*\right|^2}{2^{14} \pi^5 m_B^3} \lambda^{1/2} v \Bigg\{ \left|M_+^{+-}\right|^2 + \left|M_-^{+-}\right|^2 + \left|M_+^{++}\right|^2 + \left|M_-^{++}\right|^2 + \left|M_+^{-+}\right|^2 \\
&+ \left|M_-^{-+}\right|^2 + \left|M_+^{--}\right|^2 + \left|M_-^{--}\right|^2 + \left|M_0^{++}\right|^2 + \left|M_0^{+-}\right|^2 + \left|M_0^{-+}\right|^2 + \left|M_0^{--}\right|^2 \Bigg\}
\end{aligned} \tag{26}$$

We can now proceed to calculate the quantities $\dfrac{\Gamma_+}{\Gamma_-}$ and $\dfrac{\Gamma_L}{\Gamma_T} = \dfrac{\Gamma_0}{\Gamma_+ + \Gamma_-}$ (here, subscripts $+, -, 0$ correspond to the tensor meson helicities), lepton forward–backward asymmetry and longitudinal polarization of the final lepton. These quantities can all be measured in experiments. Since these quantities all involve "new" Wilson coefficients, they might be very sensitive to new physics.

The expressions for the quantities $\Gamma_\pm$ and $\Gamma_0$ can easily be obtained from Eq. (26), which can be written as,

$$\begin{aligned}
\Gamma_\pm &= \frac{G_F^2 \alpha^2 \left|V_{tb} V_{ts}^*\right|^2}{2^{14} \pi^5 m_B^3} \int_{4m_\ell^2}^{(m_B - m_{K_2})^2} dq^2 v \lambda^{1/2} \int_{-1}^{+1} d\cos\theta \Bigg\{ \left|M_\pm^{++}\right|^2 + \left|M_\pm^{--}\right|^2 \\
&+ \left|M_\pm^{-+}\right|^2 + \left|M_\pm^{+-}\right|^2 \Bigg\} ~, \qquad (27) \\
\Gamma_0 &= \frac{G_F^2 \alpha^2 \left|V_{tb} V_{ts}^*\right|^2}{2^{14} \pi^5 m_B^3} \int_{4m_\ell^2}^{(m_B - m_{K_2})^2} dq^2 v \lambda^{1/2} \int_{-1}^{+1} d\cos\theta \Bigg\{ \left|M_0^{++}\right|^2 + \left|M_0^{--}\right|^2 \\
&+ \left|M_0^{-+}\right|^2 + \left|M_0^{+-}\right|^2 \Bigg\} ~. \qquad (28)
\end{aligned}$$

The differential forward–backward asymmetry of the state final lepton can be obtained from Eq. (26) in the following way,

$$\frac{d\mathcal{A}_{FB}}{dq^2} = \int_0^{+1} d\cos\theta \frac{d\Gamma}{dq^2\, d\cos\theta} - \int_{-1}^0 d\cos\theta \frac{d\Gamma}{dq^2\, d\cos\theta} ~. \tag{29}$$

At the end of this section we present the expression for the longitudinal polarization of the final state lepton, which also might be very useful for establishing new physics beyond the SM. The expression for the longitudinal polarization of the final state lepton can easily be calculated from Eq. (26), which has the following form (see also [28]),

$$\mathcal{P}_L = \frac{\displaystyle\int_{4m_\ell^2}^{(m_B-m_{K_2})^2} dq^2 v\lambda^{1/2} \int_{-1}^{+1} d\cos\theta \Big[\chi_1 - \chi_2\Big]}{\displaystyle\int_{4m_\ell^2}^{(m_B-m_{K_2})^2} dq^2 v\lambda^{1/2} \int_{-1}^{+1} d\cos\theta \Big[\chi_1 + \chi_2\Big]} ~, \tag{30}$$

where

$$\begin{aligned}
\chi_1 &= \left|M_-^{-+}\right|^2 + \left|M_+^{-+}\right|^2 + \left|M_-^{--}\right|^2 + \left|M_+^{--}\right|^2 + \left|M_0^{-+}\right|^2 + \left|M_0^{--}\right|^2 ~, \\
\chi_2 &= \left|M_-^{+-}\right|^2 + \left|M_+^{+-}\right|^2 + \left|M_-^{++}\right|^2 + \left|M_+^{++}\right|^2 + \left|M_0^{+-}\right|^2 + \left|M_0^{++}\right|^2 ~.
\end{aligned}$$

# 3 Numerical analysis

In this section we investigate the dependence of the physical quantities mentioned in section 2, on the new Wilson coefficients. The main input parameters in our calculations are the new Wilson coefficients and the form factors responsible for the $B \to K_2$ transition. We use the results of [29] for the form factors, which are calculated within the QCD sum rules method. The $q^2$ dependence of all form factors are described by the following formula,

$$F_i(q^2) = \frac{F_i(0)}{1 - a_i \left(\dfrac{q^2}{m_B^2}\right) + b_i \left(\dfrac{q^2}{m_B^2}\right)^2} ~.$$

The values of parameters $F_i(0)$, $a_i$ and $b_i$ for different form factors are in Table 1 (this table is taken from [29]).

The Wilson coefficients $C_7^{eff}$ and $C_9^{eff}$ which we use in our analysis are given in Eqs. (4) and (8) with $C_9 = 4.253$ and $C_7 = -0.311$ at $\mu = m_b$ scale and $C_{10} = -4.546$.

As has already been noted, other input parameters needed are the new Wilson coefficients. A systematic analysis of the $B \to K^* \ell^+ \ell^-$ decay is carried out in [20] using the most general, model independent Hamiltonian to find the constraints on the new Wilson coefficients. In accordance with the result of [20], we will vary the vector type new Wilson coefficients in between $-C_{10}$ and $+C_{10}$. For the scalar type operators the following constraint is obtained in [20],

$$\begin{aligned}
\left|C_{LRLR}\right|^2 + \left|C_{LRRL}\right|^2 &\le 0.44 ~, \quad (\text{from } B \to \mu^+\mu^-~) ~, \\
\frac{1}{2}\left(\left|C_{LRLR}\right|^2 + \left|C_{LRRL}\right|^2\right) &\le 45 ~, \quad (\text{from } B \to X_s\mu^+\mu^-~) ~,
\end{aligned}$$

In our numerical calculations we shall use rather a broader range for the scalar type Wilson coefficients, i.e., we assume that they also vary between $-C_{10}$ and $+C_{10}$.

| | $F(0)$ | $a$ | $b$ |
|---|---|---|---|
| $A_0$ | $0.25 \pm 0.04$ | 1.57 | 0.10 |
| $A_1$ | $0.14 \pm 0.02$ | 1.21 | 0.52 |
| $A_2$ | $0.05 \pm 0.02$ | 1.32 | 14.9 |
| $V$ | $0.16 \pm 0.02$ | 2.08 | 1.50 |
| $T_1$ | $0.14 \pm 0.02$ | 2.08 | 1.50 |
| $T_2$ | $0.14 \pm 0.02$ | 1.22 | 0.35 |
| $T_3$ | $0.01^{+0.02}_{-0.01}$ | 9.91 | 276 |

Table 1: $B$ meson decay form factors in a three–parameter fit.

In Figs. (1) and (2) we present the dependence of $\Gamma_L/\Gamma_T$ on the new Wilson coefficients for $\mu$ and $\tau$ lepton channels, respectively.

We see from Fig. (1) that $\Gamma_L/\Gamma_T$ is most sensitive to the vector type coefficients, and practically insensitive to scalar type interaction. The ratio $\Gamma_L/\Gamma_T$ becomes smaller (larger) in the presence of the Wilson coefficients $C_{LL}$ and $C_{LR}$ compared to the SM value when these coefficients get negative (positive) values. This situation is to the contrary when $C_{RL}$ takes role in the numerical calculations.

For the $\tau$ channel, the ratio $\Gamma_L/\Gamma_T$ is strongly dependent on the coefficients $C_{LR}$ and $C_{RR}$, and considerable change happens also in the presence of the scalar interaction. We observe from the relevant figure that the ratio changes several times compared to the SM case.

We further study also the dependence of the asymmetry parameter $\alpha = \dfrac{\Gamma_- - \Gamma_+}{\Gamma_- + \Gamma_+}$ on the new Wilson coefficients. The result of this analysis can be summarized briefly as follows: This ratio is sensitive only to the vector type new Wilson coefficients, but insensitive to the presence of the scalar interactions. The parameter $\alpha$ exhibits strong dependence on the coefficients $C_{RL}$ and $C_{RR}$.

We present in Fig. (3) the dependence of the forward–backward asymmetry for the $\mu$ channel on $q^2$, when $C_{LL}$ is taken into account. We deduce from this figure that, there is quite a significant shift in the zero position of zero of $\mathcal{A}_{FB}$. The zero position is shifted to the right (left) for the negative (positive) values of the coefficient $C_{LL}$ compared to the SM value.

In further numerical analysis we also investigate the dependence the forward–backward asymmetry $\mathcal{A}_{FB}$ on other new Wilson coefficients for the $\mu$ and $\tau$ channels, and the outcome of these results can be summarized as follows:

In the case of $\mu$–channel

- The zero position of $\mathcal{A}_{FB}$ is shifted to the right (left) compared to the SM prediction when $C_{LR}$ gets positive (negative) values, similar to the case when $C_{LL}$ is present.

- The zero position of $\mathcal{A}_{FB}$ is insensitive to the presence of the coefficients $C_{RL}$ and $C_{RR}$, and also to all scalar interaction coefficients.

- the maximum value of $\mathcal{A}_{FB}$ is realized in the presence of the coefficient $C_{RL}$.

In the case of $\tau$–channel

- The zero position of $\mathcal{A}_{FB}$ is shifted to the right (left) when the Wilson coefficients $C_{LL}$ and $C_{LR}$ runs over negative (positive) values. The situation here is different compared to the $\mu$ case.

- The zero position of $\mathcal{A}_{FB}$ is insensitive to the presence of all remaining coefficients.

- The behavior of $\mathcal{A}_{FB}$ is very sensitive to the coefficients $C_{RL}$, $C_{LRRL}$ and $C_{RLRL}$ to the variation in $q^2$ in the range $q^2 > 13.8\ GeV^2$. It is observed that the value of $\mathcal{A}_{FB}$ is magnified 2–4 times compared to that of the SM case. Therefore the measurement of the $\mathcal{A}_{FB}$ can be very useful for establishing new physics beyond the SM.

In Figs. (4) and (5) we present the dependence of the final $\mu$ and $\tau$ leptons longitudinal polarizations on the new Wilson coefficients. We observe from these figures that,

In the case of $\mu$–channel

- $P_L$ is sensitive to the existence of all new Wilson coefficients, except the coefficients $C_{RLLR}$ and $C_{LRLR}$. The dependence of $P_L$ on $C_{LL}\,(C_{RL})$ has the tendency to increase (decrease) in the region $-C_{10} \leq C_{RL}\,(C_{LL}) \leq C_{10}$. For all other coefficients $P_L$ increase firstly in the region from $-C_{10}$ to zero (this region is from $-C_{10}$ to two for the coefficient $C_{LR}$) and then decreases when the new Wilson coefficients vary in the region from zero to $C_{10}$.

In the case of $\tau$–channel

- In this channel $P_L$ exhibits strong dependence on the coefficients $C_{LR}$, $C_{RR}$ and also on the scalar interaction coefficients $C_{LRRL}$ and $C_{RLRL}$. Therefore the measurement of the longitudinal polarization of the leptons can be quite informative about the nature and the confirmation of the new physics beyond the SM.

# 4 Conclusion

In this work the sensitivity of the physically measurable quantities, such as $\Gamma_L/\Gamma_T$, $\mathcal{A}_{FB}$ and the final lepton polarization for the $B \to K_2 \ell^+ \ell^-$ decay is investigated using the most general, model independent four–fermion interaction. It is observed that the ratio $\Gamma_L/\Gamma_T$ is quite sensitive to the new Wilson coefficients $C_{LL}$, $C_{LR}$ and $C_{RR}$ for the $B \to K_2 \mu^+ \mu^-$ channel, while for the $B \to K_2 \tau^+ \tau^-$ channel this ratio is strongly dependent on the coefficients $C_{LR}$ and $C_{RR}$. This ratio is rather weakly dependent on the scalar interaction coefficients.

We also studied in detail the dependence of the forward–backward asymmetry for both channels on $q^2$. It is found that the zero position of $\mathcal{A}_{FB}$ is shifted to right or left compared to its SM value. We also show that the value of $\mathcal{A}_{FB}$ for the $\tau$ channel is quite sensitive to the existence of scalar type interactions. The longitudinal polarization of the leptons shows sensitivity to all new Wilson coefficients, except the coefficient $C_{LRLR}$.

Measurement of these quantities can give invaluable information, not only about the existence of new physics, but also about the signs of the new Wilson coefficients.

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# Figure captions

**Fig. (1)** The dependence of the ratio of the decay widths when $K_2$ meson is longitudinally and transversally polarized on the new Wilson coefficients for the $B \to K_2 \mu^+ \mu^-$ decay.

**Fig. (2)** The same as in Fig. (1), but for the $B \to K_2 \tau^+ \tau^-$ decay.

**Fig. (3)** The dependence of the forward–backward asymmetry on $q^2$ at several fixed values of the Wilson coefficient $C_{LL}$ for the $B \to K_2 \mu^+ \mu^-$ decay.

**Fig. (4)** The dependence of the longitudinal lepton polarization of the $\mu$–lepton on the new Wilson coefficients.

**Fig. (5)** The same as in Fig. (4), but for the $\tau$–lepton.

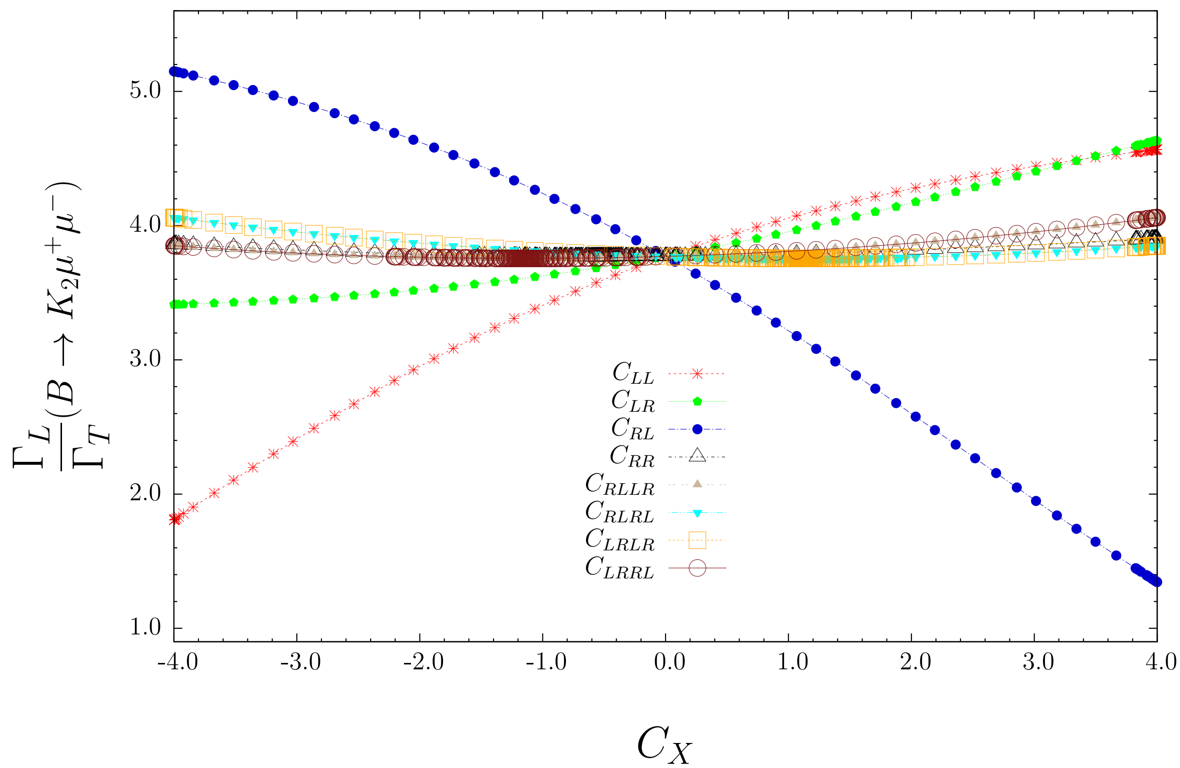


Figure 1:

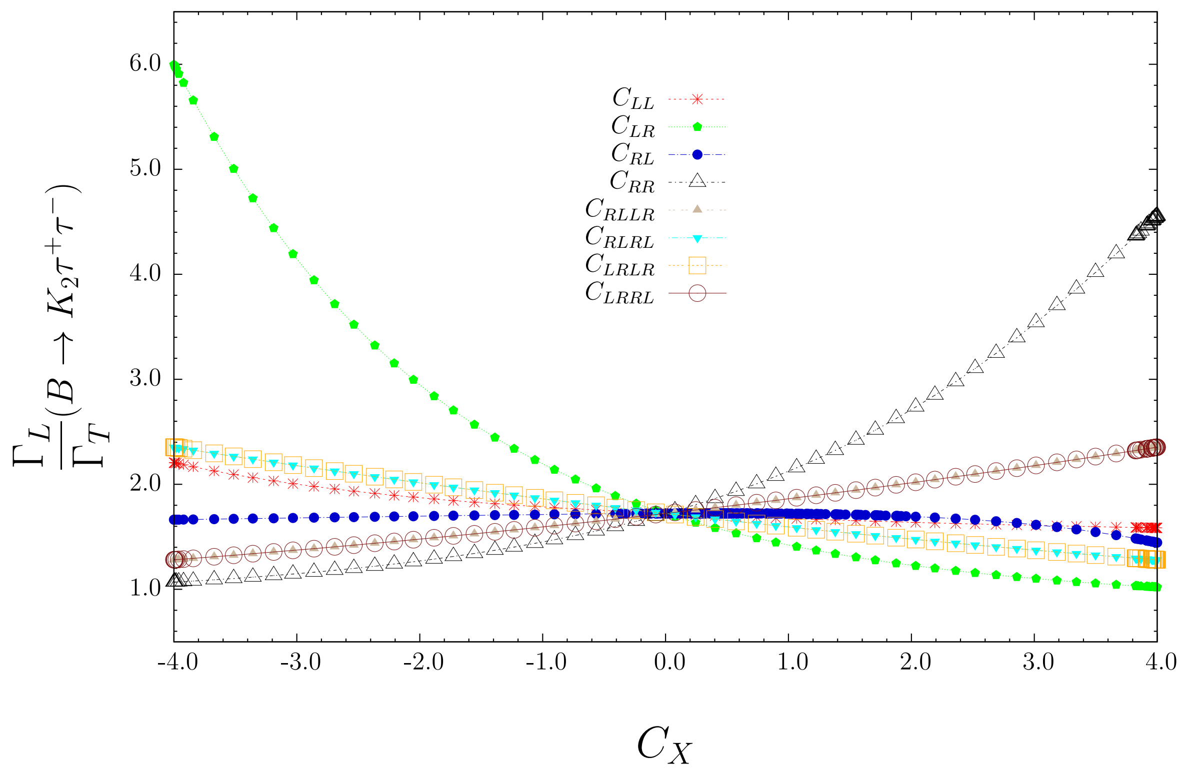


Figure 2:

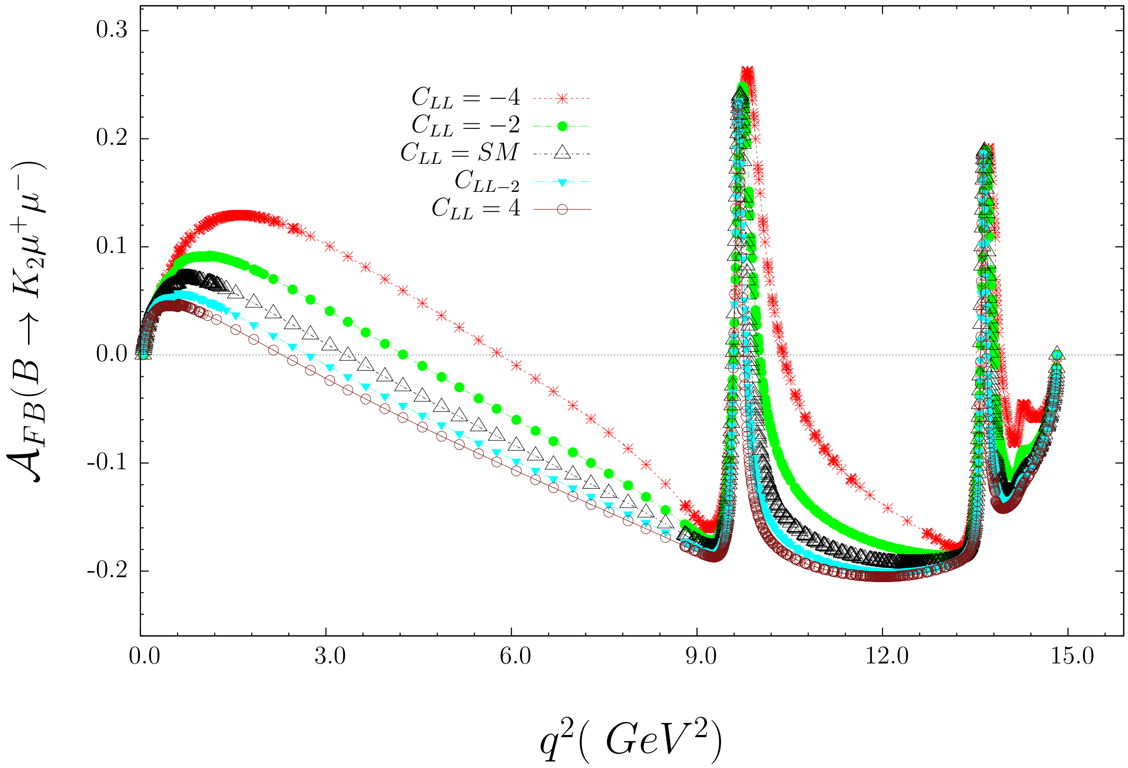


Figure 3:

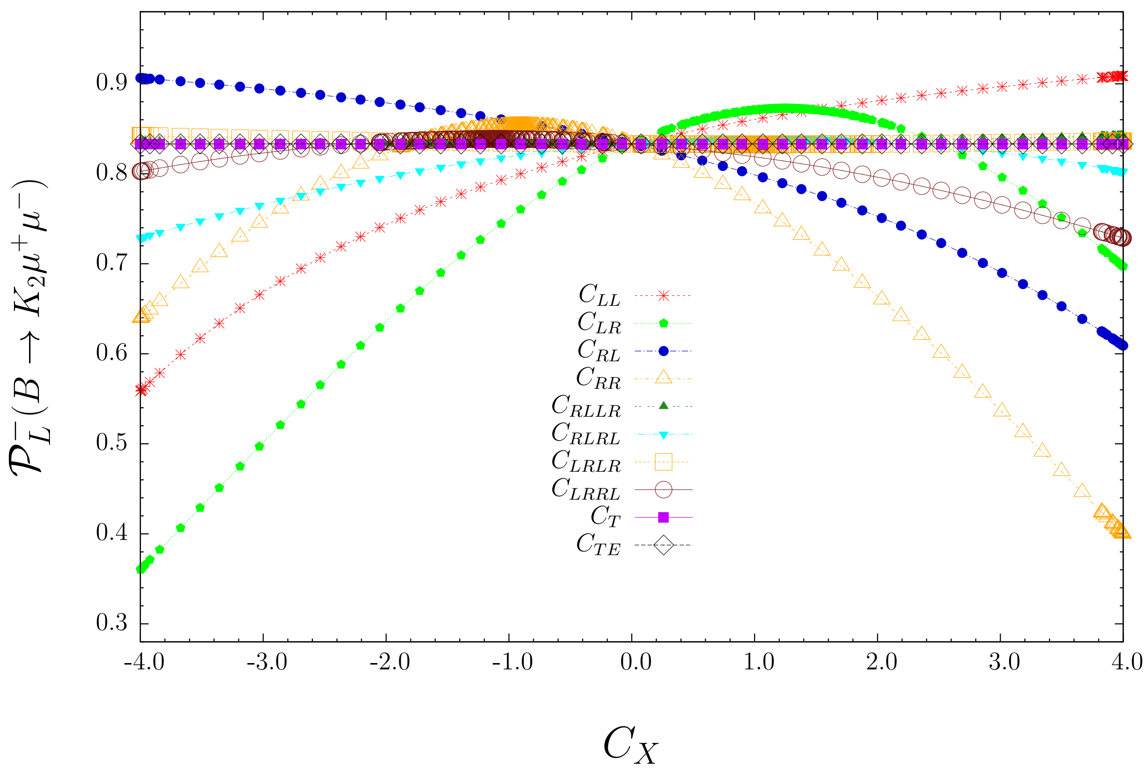


Figure 4:

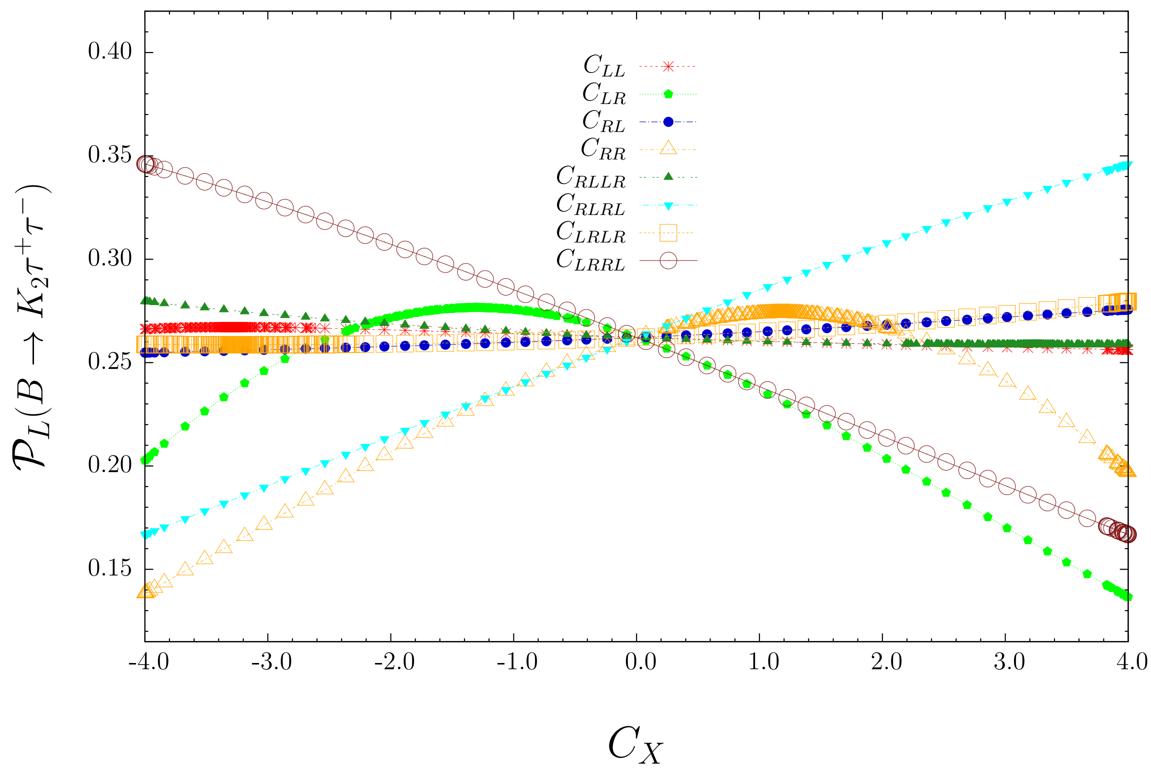
$C_{LL}$
$C_{LR}$
$C_{RL}$
$C_{RR}$
$C_{RLLR}$
$C_{RLRL}$
$C_{LRLR}$
$C_{LRRL}$
$\mathcal{P}_L(B \to K_2 \tau^+ \tau^-)$
$C_X$
0.40
0.35
0.30
0.25
0.20
0.15
-4.0
-3.0
-2.0
-1.0
0.0
1.0
2.0
3.0
4.0

Figure 5: