Exclusive $B \to K^* \ell^+ \ell^-$ decay with polarized K^* and new physics effects

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Abstract

Using the most general, model independent effective Hamiltonian, the branching ratio of the $B \to K^* \ell^+ \ell^-$ decay, when K^* meson is longitudinally or transversally polarized, is presented. The dependence of the branching ratio on the new Wilson coefficients, when K^* meson is polarized, is studied. It is observed that the branching ratio is very sensitive to the vector and tensor type interactions, which will be useful in search of new physics beyond the Standard Model.

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1 Introduction

Rare B meson decays, induced by flavor–changing neutral current (FCNC) $b \to s(d)$ transitions, provide potentially stringiest tests of the Standard Model (SM) in flavor sector. These transitions take place in the SM at loop level, is very sensitive to the gauge structure of the SM. Moreover $b \to s(d)\ell^+\ell^-$ decay is known to be very sensitive to the various extensions of the SM. New physics effects manifest themselves in rare B meson decays in two different ways, either through new contributions to the Wilson coefficients existing in the SM or through the new structures in the effective Hamiltonian which are absent in the SM. Note that $b \to s(d)\ell^+\ell^-$ transition has been extensively studied in framework of the SM and its various extensions [1]–[15].

The rare inclusive decays are theoretically much cleaner than the exclusive decays, which require the knowledge of form factors, are also more difficult to measure. However, FCNC exclusive semileptonic decays, in particular $B \to K^*(K)\ell^+\ell^-$ will be measured precisely in the future experiments at CLEO and B–factories LHC, HERA etc. As has already been noted, new physics effects in the rare B meson decays can appear in two different ways, either through new contributions to the existing in the in the SM through the new operators in the effective Hamiltonian which are absent in the SM. Using these approaches, the $B \to K^*\ell^+\ell^-$ decay was studied in [16, 17] using the most the most general form of effective Hamiltonian that includes all possible form of interactions. It was shown that different physical observables like branching ratio, forward–backward asymmetry etc. are very sensitive to the new Wilson coefficients. One efficient way in establishing new physics effects beyond the SM is taking into account polarization effects. Along these lines these effects have been studied for the $B \to K^*\ell^+\ell^-$ decay in [15], [18]–[26]. It is shown in [26] that there exists region of new Wilson coefficients in which the decay rate agrees with the SM prediction while lepton polarization does not. In other words, in this region of new Wilson coefficients new physics effects can be established by measuring lepton polarization only but not the branching ratio.

In this connection there follows the following question. How sensitive is the branching ratio to the new Wilson coefficients when K^* meson is polarized longitudinally or transversally? The goal of the present work is to find an answer to this question.

The paper is organized as follows. In section 2, using a general form of four–Fermi interaction we derive the model independent expressions for the longitudinal, transversal and normal polarizations of leptons. In section 3 we investigate the dependence of the branching ratios on the four–Fermi interactions when K^* meson is polarized transversally or longitudinally.

2 Theoretical background

In this section we calculate the branching ratio of the $B \to K^*\ell^+\ell^-$ decay when K^* meson is polarized transversally or longitudinally, using the most general, model independent four–Fermi interactions. The effective Hamiltonian for the $b \to s\ell^+\ell^-$ transition in terms of twelve model independent four–Fermi interactions can be written as [19]

$$
\begin{aligned}
\mathcal{H}_{eff} &= \frac{G\alpha}{\sqrt{2}\pi} V_{ts}V_{tb}^{*} \Bigg\{ C_{SL}\, \bar{s} i \sigma_{\mu\nu} \frac{q^{\nu}}{q^2}\, L\, b\, \bar{\ell}\gamma^{\mu}\ell + C_{BR}\, \bar{s} i \sigma_{\mu\nu} \frac{q^{\nu}}{q^2}\, R\, b\, \bar{\ell}\gamma^{\mu}\ell \\
&+ C_{LL}^{tot}\, \bar{s}_L \gamma_\mu b_L\, \bar{\ell}_L \gamma^{\mu} \ell_L + C_{LR}^{tot}\, \bar{s}_L \gamma_\mu b_L\, \bar{\ell}_R \gamma^{\mu} \ell_R + C_{RL}\, \bar{s}_R \gamma_\mu b_R\, \bar{\ell}_L \gamma^{\mu} \ell_L \\
&+ C_{RR}\, \bar{s}_R \gamma_\mu b_R\, \bar{\ell}_R \gamma^{\mu} \ell_R + C_{LRLR}\, \bar{s}_L b_R\, \bar{\ell}_L \ell_R + C_{RLLR}\, \bar{s}_R b_L\, \bar{\ell}_L \ell_R \\
&+ C_{LRRL}\, \bar{s}_L b_R\, \bar{\ell}_R \ell_L + C_{RLRL}\, \bar{s}_R b_L\, \bar{\ell}_R \ell_L + C_T\, \bar{s}\sigma_{\mu\nu} b\, \bar{\ell}\sigma^{\mu\nu}\ell \\
&+ i C_{TE}\, \epsilon^{\mu\nu\alpha\beta} \bar{s}\sigma_{\mu\nu} b\, \bar{\ell}\sigma_{\alpha\beta}\ell \Bigg\} ,
\end{aligned} \tag{1}
$$

where the chiral projection operators L and R in (1) are defined as

$$L = \frac{1-\gamma_5}{2} , \qquad R = \frac{1+\gamma_5}{2} ,$$

and C_X are the coefficients of the four–Fermi interactions, and part of these coefficients exist in the SM as well. The first two of these coefficients, C_{SL} and C_{BR}, are the nonlocal Fermi interactions which correspond to $-2m_s C_7^{eff}$ and $-2m_b C_7^{eff}$ in the SM, respectively. The following four terms in this expression are the vector type interactions with coefficients C_{LL}, C_{LR}, C_{RL} and C_{RR}. Two of these vector interactions containing C_{LL}^{tot} and C_{LR}^{tot} do also exist in the SM in combinations of the form $(C_9^{eff} - C_{10})$ and $(C_9^{eff} + C_{10})$, respectively. Therefore one can say that C_{LL}^{tot} and C_{LR}^{tot} represent sum the contributions from SM and the new physics, whose explicit forms can be written as

$$
\begin{aligned}
C_{LL}^{tot} &= C_9^{eff} - C_{10} + C_{LL} , \\
C_{LR}^{tot} &= C_9^{eff} + C_{10} + C_{LR} .
\end{aligned}
$$

The terms with coefficients C_{LRLR}, C_{RLLR}, C_{LRRL} and C_{RLRL} describe the scalar type interactions. The remaining two terms with the coefficients C_T and C_{TE}, obviously, describe the tensor type interactions.

Exclusive decay $B \to K^* \ell^+ \ell^-$ is described in terms of matrix elements of the quark operators over meson states, which are parametrized in terms of form factors. It follows from Eq. (1) that, in order to calculate the amplitude of the $B \to K^* \ell^+ \ell^-$ decay the following matrix elements are needed

$$
\begin{aligned}
&\left\langle K^* \left| \bar{s}\gamma_\mu (1 \pm \gamma_5) b \right| B \right\rangle , \\
&\left\langle K^* \left| \bar{s} i \sigma_{\mu\nu} q^{\nu} (1 \pm \gamma_5) b \right| B \right\rangle , \\
&\left\langle K^* \left| \bar{s} (1 \pm \gamma_5) b \right| B \right\rangle , \\
&\left\langle K^* \left| \bar{s}\sigma_{\mu\nu} b \right| B \right\rangle .
\end{aligned}
$$

These matrix elements are defined as follows:

$$
\begin{aligned}
&\left\langle K^*(p_{K^*}, \varepsilon) \left| \bar{s}\gamma_\mu (1 \pm \gamma_5) b \right| B(p_B) \right\rangle = \\
&\quad -\epsilon_{\mu\nu\lambda\sigma} \varepsilon^{*\nu} p_{K^*}^{\lambda} q^{\sigma} \frac{2V(q^2)}{m_B + m_{K^*}} \pm i \varepsilon_\mu^* (m_B + m_{K^*}) A_1(q^2)
\end{aligned} \tag{2}
$$

$$\mp i(p_B+p_{K^*})_\mu(\varepsilon^* q)\frac{A_2(q^2)}{m_B+m_{K^*}} \mp i q_\mu \frac{2m_{K^*}}{q^2}(\varepsilon^* q)\Big[A_3(q^2)-A_0(q^2)\Big] ~,$$

$$\begin{aligned}
&\left\langle K^*(p_{K^*},\varepsilon)\left|\bar s i\sigma_{\mu\nu}q^\nu(1\pm\gamma_5)b\right|B(p_B)\right\rangle = \\
&\quad 4\epsilon_{\mu\nu\lambda\sigma}\varepsilon^{*\nu}p_{K^*}^\lambda q^\sigma T_1(q^2) \pm 2i\Big[\varepsilon_\mu^*(m_B^2-m_{K^*}^2)-(p_B+p_{K^*})_\mu(\varepsilon^* q)\Big]T_2(q^2) \\
&\quad \pm 2i(\varepsilon^* q)\left[q_\mu-(p_B+p_{K^*})_\mu\frac{q^2}{m_B^2-m_{K^*}^2}\right]T_3(q^2) ~,
\end{aligned} \tag{3}$$

$$\begin{aligned}
&\left\langle K^*(p_{K^*},\varepsilon)\left|\bar s \sigma_{\mu\nu} b\right|B(p_B)\right\rangle = \\
&\quad i\epsilon_{\mu\nu\lambda\sigma}\Bigg\{-2T_1(q^2)\varepsilon^{*\lambda}(p_B+p_{K^*})^\sigma+\frac{2}{q^2}(m_B^2-m_{K^*}^2)\Big[T_1(q^2)-T_2(q^2)\Big]\varepsilon^{*\lambda}q^\sigma \\
&\quad -\frac{4}{q^2}\left[T_1(q^2)-T_2(q^2)-\frac{q^2}{m_B^2-m_{K^*}^2}T_3(q^2)\right](\varepsilon^* q)p_{K^*}^\lambda q^\sigma\Bigg\} ~.
\end{aligned} \tag{4}$$

where $q=p_B-p_{K^*}$ is the momentum transfer and ε is the polarization vector of K^* meson. In order to ensure finiteness of (2) and (4) at $q^2=0$, we assume that $A_3(q^2=0)=A_0(q^2=0)$ and $T_1(q^2=0)=T_2(q^2=0)$. The matrix element $\left\langle K^*\left|\bar s(1\pm\gamma_5)b\right|B\right\rangle$ can be calculated from Eq. (2) by contracting both sides of Eq. (2) with q^μ and using equation of motion. Neglecting the mass of the strange quark in this matrix element, we get element

$$\left\langle K^*(p_{K^*},\varepsilon)\left|\bar s(1\pm\gamma_5)b\right|B(p_B)\right\rangle = \frac{1}{m_b}\Big[\mp 2im_{K^*}(\varepsilon^* q)A_0(q^2)\Big] ~. \tag{5}$$

In deriving Eq. (5) we have used the exact relation

$$2m_{K^*}A_3(q^2)=(m_B+m_{K^*})A_1(q^2)-(m_B-m_{K^*})A_2(q^2) ~.$$

Taking into account Eqs. (1–5), the matrix element of the $B\to K^*\ell^+\ell^-$ decay can be written as

$$\begin{aligned}
&\mathcal{M}(B\to K^*\ell^+\ell^-) = \frac{G\alpha}{4\sqrt{2}\pi}V_{tb}V_{ts}^* \\
&\quad\times\Bigg\{\bar\ell\gamma^\mu\ell\Big[-2A\epsilon_{\mu\nu\lambda\sigma}\varepsilon^{*\nu}p_{K^*}^\lambda q^\sigma-iB\varepsilon_\mu^*+iC(\varepsilon^* q)(p_B+p_{K^*})_\mu+iD(\varepsilon^* q)q_\mu\Big] \\
&\quad+\bar\ell\gamma^\mu\gamma_5\ell\Big[-2E\epsilon_{\mu\nu\lambda\sigma}\varepsilon^{*\nu}p_{K^*}^\lambda q^\sigma-iF\varepsilon_\mu^*+iG(\varepsilon^* q)(p_B+p_{K^*})_\mu+iH(\varepsilon^* q)q_\mu\Big] \\
&\quad+\bar\ell\ell\Big[iQ(\varepsilon^* q)\Big]+\bar\ell\gamma_5\ell\Big[iN(\varepsilon^* q)\Big] \\
&\quad+4\bar\ell\sigma^{\mu\nu}\ell\Big(iC_T\epsilon_{\mu\nu\lambda\sigma}\Big)\Big[-2T_1\varepsilon^{*\lambda}(p_B+p_{K^*})^\sigma+B_6\varepsilon^{*\lambda}q^\sigma-B_7(\varepsilon^* q){p_{K^*}}^\lambda q^\sigma\Big] \\
&\quad+16C_{TE}\bar\ell\sigma_{\mu\nu}\ell\Big[-2T_1\varepsilon^{*\mu}(p_B+p_{K^*})^\nu+B_6\varepsilon^{*\mu}q^\nu-B_7(\varepsilon^* q){p_{K^*}}^\mu q^\nu\Big]\Bigg\} ~.
\end{aligned} \tag{6}$$

The auxiliary functions in Eq. (6) are given by

$$A = (C_{LL}^{tot}+C_{LR}^{tot}+C_{RL}+C_{RR})\frac{V}{m_B+m_{K^*}}-4(C_{BR}+C_{SL})\frac{T_1}{q^2} ~,$$

$$
\begin{aligned}
B &= (C_{LL}^{tot} + C_{LR}^{tot} - C_{RL} - C_{RR})(m_B + m_{K^*})A_1 - 4(C_{BR} - C_{SL})(m_B^2 - m_{K^*}^2)\frac{T_2}{q^2} ~, \\
C &= (C_{LL}^{tot} + C_{LR}^{tot} - C_{RL} - C_{RR})\frac{A_2}{m_B + m_{K^*}} - 4(C_{BR} - C_{SL})\frac{1}{q^2}\left[T_2 + \frac{q^2}{m_B^2 - m_{K^*}^2}T_3\right] ~, \\
D &= 2(C_{LL}^{tot} + C_{LR}^{tot} - C_{RL} - C_{RR})m_{K^*}\frac{A_3 - A_0}{q^2} + 4(C_{BR} - C_{SL})\frac{T_3}{q^2} ~, \\
E &= (C_{LR}^{tot} + C_{RR} - C_{LL}^{tot} - C_{RL})\frac{V}{m_B + m_{K^*}} ~, \\
F &= (C_{LR}^{tot} - C_{RR} - C_{LL}^{tot} + C_{RL})(m_B + m_{K^*})A_1 ~, \\
G &= (C_{LR}^{tot} - C_{RR} - C_{LL}^{tot} + C_{RL})\frac{A_2}{m_B + m_{K^*}} ~, \\
H &= 2(C_{LR}^{tot} - C_{RR} - C_{LL}^{tot} + C_{RL})m_{K^*}\frac{A_3 - A_0}{q^2} ~, \\
Q &= -2(C_{LRRL} - C_{RLRL} + C_{LRLR} - C_{RLLR})\frac{m_{K^*}}{m_b}A_0 ~, \\
N &= -2(C_{LRLR} - C_{RLLR} - C_{LRRL} + C_{RLRL})\frac{m_{K^*}}{m_b}A_0 ~, \\
B_6 &= 2(m_B^2 - m_{K^*}^2)\frac{T_1 - T_2}{q^2} ~, \\
B_7 &= \frac{4}{q^2}\left(T_1 - T_2 - \frac{q^2}{m_B^2 - m_{K^*}^2}T_3\right) ~.
\end{aligned} \tag{7}
$$

The form of Eq. (6) reflects the fact that its difference from the SM case is due to the last four structures, namely, scalar and tensor type interactions. The next task to be considered is calculation of the branching ratio of the $B \to K^* \ell^+ \ell^-$ decay, when K^* is polarized transversally or longitudinally. From matrix element (6)it is easy to derive the invariant dilepton mass spectrum for the $B \to K^* \ell^+ \ell^-$ decay corresponding to the transversally and longitudinally polarized K^* meson:

$$
\frac{d\Gamma_\pm}{ds} = \frac{G^2 \alpha^2 \left|V_{tb} V_{ts}^*\right|^2}{2^{14} \pi^5} m_B \sqrt{\lambda(1,r,s)}\, v \Delta_\pm ~, \tag{8}
$$

where

$$
\begin{aligned}
\Delta_\pm &= 256 m_B^2 m_\ell \,\mathrm{Re}\Big[\Big(\sqrt{\lambda}\, m_B^2 A^* \mp B^*\Big) \Big(\sqrt{\lambda}\, C_T T_1 \pm 2 C_{TE} T_1 (1-r) \mp s B_6 C_{TE}\Big) \Big] \\
&+ \frac{4}{3} m_B^2 s \Big[\Big(3 - v^2\Big) \Big| B \mp \sqrt{\lambda}\, m_B^2 A \Big|^2 + 2 v^2 \Big| F \mp \sqrt{\lambda}\, m_B^2 E \Big|^2 \Big] \\
&+ \frac{256}{3} m_B^4 \Big[v^2 \left|C_T\right|^2 + 4 \Big(3 - 2v^2\Big) \left|C_{TE}\right|^2 \Big] \left| s B_6 - 2(1-r) T_1 \right|^2 \\
&+ \frac{1024}{3} \lambda m_B^4 \left|T_1\right|^2 \Big[\Big(3 - 2v^2\Big) \left|C_T\right|^2 + 4 v^2 \left|C_{TE}\right|^2 \Big] \\
&\pm \frac{2048}{3} \sqrt{\lambda}\, m_B^4 \Big\{ 2(1-r) \Big(3 - v^2\Big) \left|T_1\right|^2 \,\mathrm{Re}(C_T C_{TE}^*) \\
&- s \,\mathrm{Re}\Big(\Big[v^2 C_T^* C_{TE} + (3 - 2v^2) C_T C_{TE}^* \Big] B_6^* T_1 \Big) \Big\} ~,
\end{aligned} \tag{9}
$$

and,

$$\frac{d\Gamma_\pm}{ds} = \frac{G^2\alpha^2 \left|V_{tb}V_{ts}^*\right|^2}{2^{14}\pi^5} m_B \sqrt{\lambda(1,r,s)}\, v\Delta_0 \ , \tag{10}$$

where,

$$\begin{aligned}
\Delta_0 &= \frac{4}{r}\lambda m_B^2 m_\ell \,\mathrm{Re}\Big(- F + m_B^2(1-r)G + m_B^2 sH\Big)N^* \\
&+ \frac{1}{r}\lambda m_B^4 \Big\{ sv^2 \left|Q\right|^2 + \frac{1}{3}\lambda m_B^2 (3-v^2) \left|C\right|^2 - \frac{2}{3}(1-r-s)(3-v^2)\,\mathrm{Re}(BC^*) \\
&- \frac{2}{3}\Big[(1-r-s)(3-v^2) + 3s(1-v^2)\Big]\,\mathrm{Re}(FG^*) - 2s(1-v^2)\,\mathrm{Re}(FH^*) \\
&+ s\left|N\right|^2 + m_B^2 s^2(1-v^2)\left|H\right|^2 + 2m_B^2 s(1-r)(1-v^2)\,\mathrm{Re}(GH^*)\Big\} \\
&+ \frac{1}{3r}m_B^2\Big\{(\lambda+4rs)(3-v^2)\left|B\right|^2 + \lambda m_B^4\Big[\lambda(3-v^2) - 3s(s-2r-2)(1-v^2)\Big]\left|G\right|^2 \\
&+ \Big[\lambda(3-v^2) + 8rsv^2\Big]\left|F\right|^2 \Big\} \\
&+ \frac{64}{r}m_B^2 m_\ell\,\mathrm{Re}\Big(B_6 C_{TE}\Big[(\lambda+4rs)B^* - \lambda m_B^2(1-r-s)C^*\Big]\Big) \\
&+ \frac{32}{r}\lambda m_B^4 m_\ell\,\mathrm{Re}\Big(B_7 C_{TE}\Big[\lambda m_B^2 C^* - (1-r-s)B^*\Big]\Big) \\
&+ \frac{16}{3r}m_B^4 s\Big\{\lambda^2 m_B^4\left|B_7\right|^2 + 4(\lambda+4rs)\left|B_6\right|^2 - 4\lambda m_B^2(1-r-s)\,\mathrm{Re}(B_6B_7^*) \\
&- 16[\lambda+4r(1-r)]\,\mathrm{Re}(B_6T_1^*) + 8\lambda m_B^2(1+3r-s)\,\mathrm{Re}(B_7T_1^*) \\
&+ 16(1+3r-s)^2\left|T_1\right|^2\Big\} \times \Big\{v^2\left|C_T\right|^2 + 4(3-2v^2)\left|C_{TE}\right|^2\Big\} \\
&- \frac{128}{r}m_B^2 m_\ell\Big\{\Big[\lambda+4r(1-r)\Big]\,\mathrm{Re}(C_{TE}T_1B^*) - \lambda m_B^2(1+3r-s)\,\mathrm{Re}(C_{TE}T_1C^*)\Big\} \ .
\end{aligned} \tag{11}$$

In Eqs. (8) and (10) subscripts $\pm$ and 0 denote polarization of the K^* meson, $v = \sqrt{1-4m_\ell^2/(m_B^2 s)}$ is the lepton velocity, $\lambda(1,r,s) = 1+r^2+s^2-2r-2s-2rs$, $r = m_{K^*}^2/m_B^2$ and $s = q^2/m_B^2$.

3 Numerical analysis

We first present the main input parameters which have been used in the present work whose values are: $|V_{tb}V_{ts}^*| = 0.0385$, $\alpha^{-1} = 129$, $G_F = 1.17\times 10^{-5}\ GeV^{-2}$, $\Gamma_B = 4.22\times 10^{-13}\ GeV$, $C_9^{eff} = 4.344$, $C_{10} = -4.669$. This value of the Wilson coefficient C_9^{eff} corresponds only to short distance contribution. In addition to the short distance contribution, it is well known that C_9^{eff} also receives long distance contributions associated with the real $\bar{c}c$ intermediate states, i.e., with the J/ψ family. In this work we restricted ourselves only to short distance contributions. As far as C_7^{eff} is concerned, experimental results fixes only the modulo of it. For this reason throughout our analysis we have considered both possibilities, i.e., $C_7^{eff} = \mp 0.313$, where the upper sign corresponds to the SM prediction. The values of the input parameters which are summarized above, have been fixed by their central values.

In performing numerical calculations we also need the explicit form of the form factors and for this purpose we have used the results of [27] (see also [28]) where the radiative corrections to the leading twist contribution and $SU(3)$ breaking effects are also taken into account. In this work the q^2 dependence of the form factors are given in terms of three parameters as

$$F(q^2) = \frac{F(0)}{1 - a_F \dfrac{q^2}{m_B^2} + b_F \left(\dfrac{q^2}{m_B^2}\right)^2} ~,$$

where the values of parameters $F(0)$, a_F and b_F for the $B \to K^*$ decay are listed in Table 1.

	$F(0)$	a_F	b_F
$A_1^{B\to K^*}$	0.34 ± 0.05	0.60	-0.023
$A_2^{B\to K^*}$	0.28 ± 0.04	1.18	0.281
$V^{B\to K^*}$	0.46 ± 0.07	1.55	0.575
$T_1^{B\to K^*}$	0.19 ± 0.03	1.59	0.615
$T_2^{B\to K^*}$	0.19 ± 0.03	0.49	-0.241
$T_3^{B\to K^*}$	0.13 ± 0.02	1.20	0.098

Table 1: B meson decay form factors in a three-parameter fit, where the radiative corrections to the leading twist contribution and SU(3) breaking effects are taken into account.

We present our numerical results in a series of graphs. In Figs. (1) and (2) the dependence of the branching ratio $\mathcal{B}^{\pm}(B \to K^* \mu^+ \mu^-)$ on the new Wilson coefficients is depicted, where superscripts $\pm$ correspond to the polarization of K^* meson. From these figures we observe that the branching ratio in both cases depends quite strongly on the on tensor interaction. From Fig. (1) we see that branching ratio $\mathcal{B}^+$ is sensitive to the vector interactions with coefficients C_{RR} and C_{RL}, while $\mathcal{B}^-$ is more sensitive to the coefficient C_{LL}. It further follows from these figures that terms proportional to C_{RL}, C_{RR} give constructive contribution to the branching ratios $\mathcal{B}^+$ and the ones proportional to C_{LL}, C_{LR} do so to the branching ratio $\mathcal{B}^-$, respectively. For both cases the contribution coming from scalar part is quite small. From these figures we also see that $\mathcal{B}^- > \mathcal{B}^+$. This fact can easily be understood from Eq. (10). In SM, in the limit $m_\ell \to 0$ we get

$$\begin{aligned}
\Delta_\pm &= \left| 2C_9 m_B \left[(1+r)A_1 \mp \sqrt{\lambda} \frac{V}{1+r} \right] + 8C_7 \frac{m_b}{s} \left(T_2 \mp \sqrt{\lambda} T_1 \right) \right|^2 \\
&+ 4m_B^2 \left| C_{10} \right|^2 \left| (1+r)A_1 \mp \sqrt{\lambda} \frac{V}{1+r} \right|^2 ~,
\end{aligned}$$

from which it obviously follows that $\mathcal{B}^- > \mathcal{B}^+$.

For the $B \to K^* \tau^+ \tau^-$ decay, apart from the magnitudes of the branching ratios $\mathcal{B}^+$ and $\mathcal{B}^-$, which become smaller compared to the muon case, all results obtained for the $B \to K^* \mu^+ \mu^-$ case remain valid for the τ lepton case as well; i.e., $\mathcal{B}^+$ and $\mathcal{B}^-$ depend strongly on C_{RR}, C_{RL} and C_{LL}, C_{LR}, respectively.

In Fig. (3) the dependence of the $B \to K^* \mu^+ \mu^-$ decay on the new Wilson coefficients when K^* meson is polarized longitudinally, is studied. This figure depicts that the branching ratio is very sensitive to all type of vector and tensor interactions. Note that, for simplicity all new Wilson coefficients in this work are assumed to be real and varied in the region between -4 and +4. From this figure we see that when C_{LL} and C_{RL} increases from -4 to +4, the branching ratios increases and decreases, respectively, while the dependence of the branching ratio on the tensor interactions have a rather symmetrical form on both sides of the origin. In other words up to zero values of the tensor interaction coefficients the branching ratio decreases (for the $B \to K^* \tau^+ \tau^-$ case this symmetry point is slightly shifted) and it increases from 0 to +4.

Depicted in Fig. (4) is the dependence of the branching ratio on the new Wilson coefficients when K^* is transversally polarized ($\mathcal{B}_T = \mathcal{B}^+ + \mathcal{B}^-$). Obviously this figure for muon decay channel demonstrates strong dependence on tensor interaction coefficients and on the coefficient C_{LL}. In the case of $B \to K^* \tau^+ \tau^-$ decay this branching ratio is strongly dependent on tensor interactions.

Finally, in Figs. (5) and (6) we present the dependence of another physically measurable quantity, namely the ratio of the branching ratios $\mathcal{B}_L/\mathcal{B}_T$ on new Wilson coefficients. From these figures we conclude that dominant contribution for the $B \to K^* \mu^+ \mu^-$ decay comes from C_{RL}. So if in future experiments a larger value for this ratio is observed than the SM prediction, this result can be attributed solely to the vector interaction with coefficient C_{RL}, whose range is $-4 \le C_{RL} \le 0$. However if in these experiments smaller values for the same ratio is measured, this departure from the SM prediction can be explained with the help of different mechanisms. Experimentally measured value can give us information which mechanism is responsible for such a discrepancy. For the $B \to K^* \tau^+ \tau^-$ case, a measurement of the same ratio which yields $\mathcal{B}_L/\mathcal{B}_T > 1.3 \times (\mathcal{B}_L/\mathcal{B}_T)_{SM}$ indicates the existence of new vector type interaction with coefficient C_{LR}. It should be noted here that all these calculations are performed for the choice of $C_7^{eff} = -0.313$. Numerical analysis for this choice shows that all conclusions which have been made for the $C_7^{eff} = +0.313$ case remains valid, apart from a slight decrease in the magnitude of the branching ratios.

It follows from all these discussions that, the branching ratios when K^* meson has $\pm$ and zero helicities and the ratio of the branching ratios when K^* meson is polarized longitudinally and transversally, are very sensitive to the presence of different new Wilson coefficients. Experimental measurement of the branching ratio for the $B \to K^* \mu^+ \mu^-$ and $B \to K^* \tau^+ \tau^-$ decays when K^* meson has different polarizations, can give quite valuable information about new physics.

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Figure captions

Fig. (1) The dependence of the branching ratio $\mathcal{B}^+$ on the new Wilson coefficients for the $B \to K^* \mu^- \mu^+$ decay. The superscript $+$ corresponds to the polarization of K^* meson.

Fig. (2) The same as in Fig. (1), but for the $-$ polarization of K^* meson.

Fig. (3) The dependence of the branching ratio $\mathcal{B}_L$ on the new Wilson coefficients for the $B \to K^* \mu^- \mu^+$ decay. The subscript L denotes the longitudinal polarization of the K^* meson.

Fig. (4) The same as in Fig. (3), but for the case when K^* meson is polarized transversally.

Fig. (5) The dependence of the ratio of the branching ratios $\mathcal{B}_L/\mathcal{B}_T$ on new Wilson coefficients for the $B \to K^* \mu^+ \mu^-$ decay.

Fig. (6) The same as in Fig. (5), but for the $B \to K^* \tau^- \tau^+$ decay.

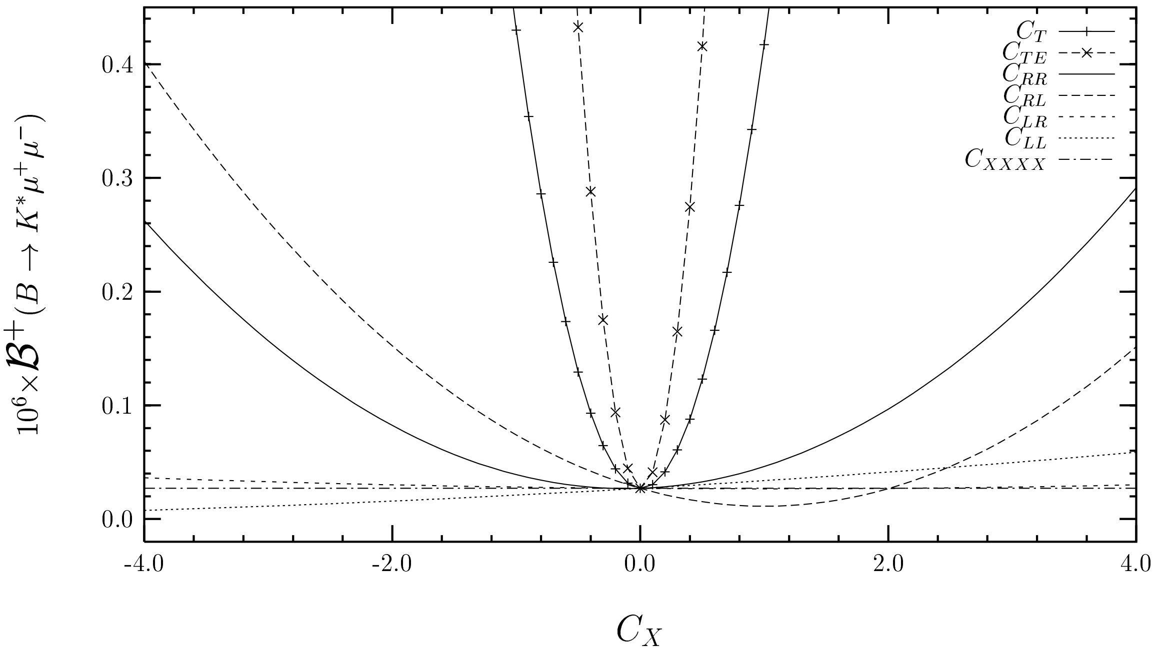
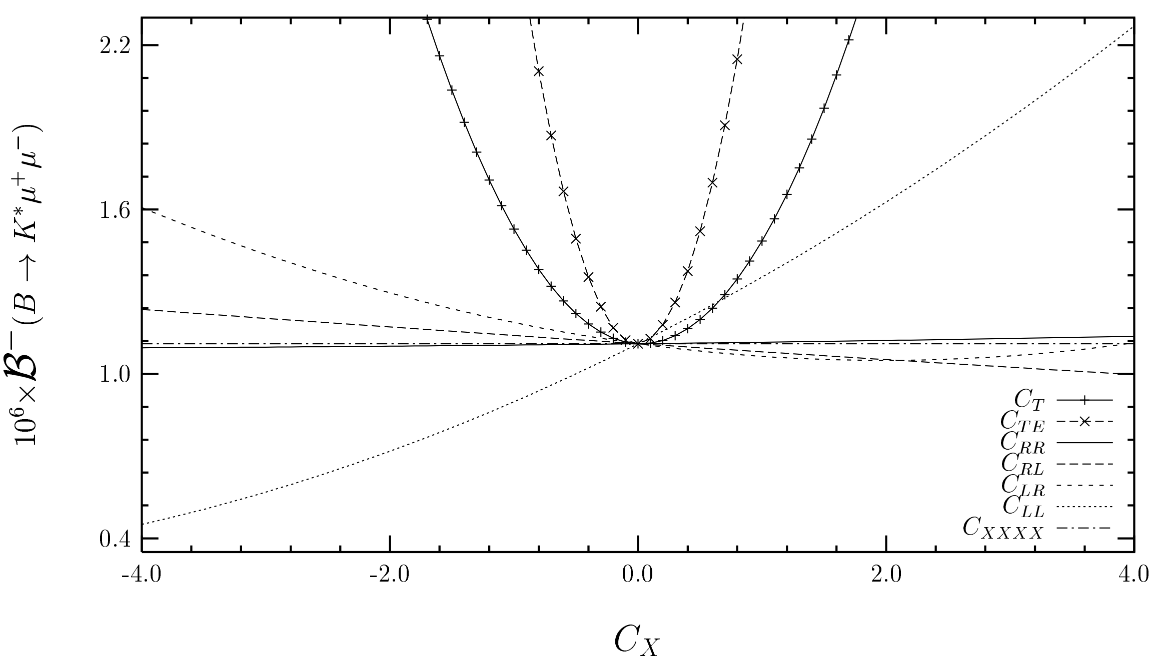


Figure 1:



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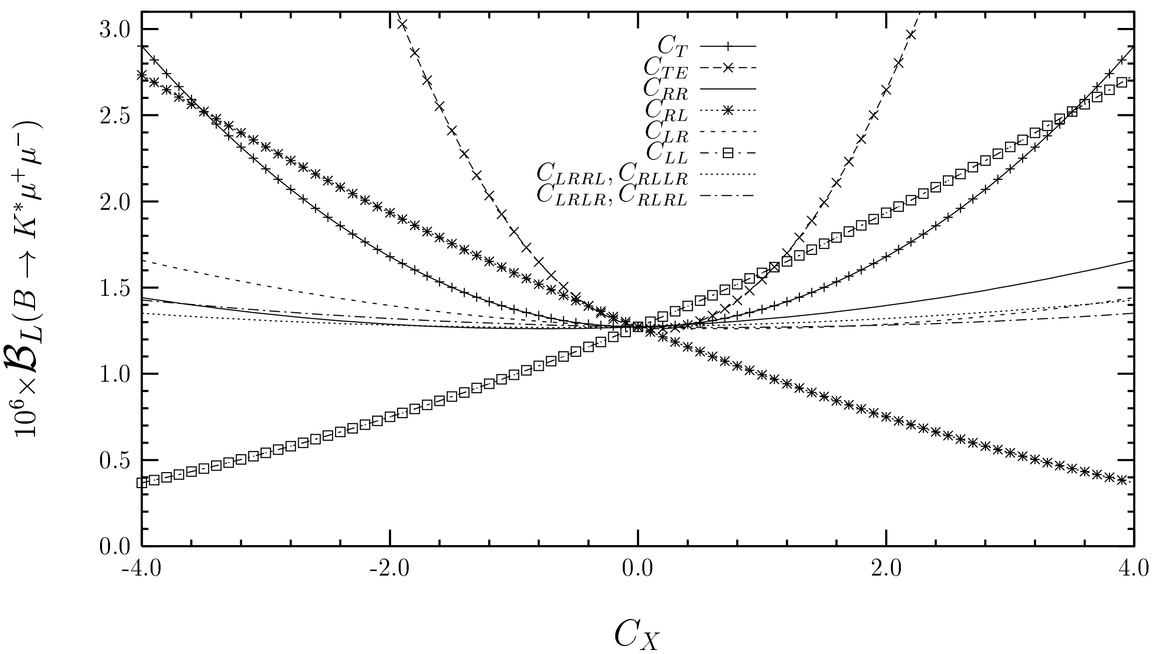

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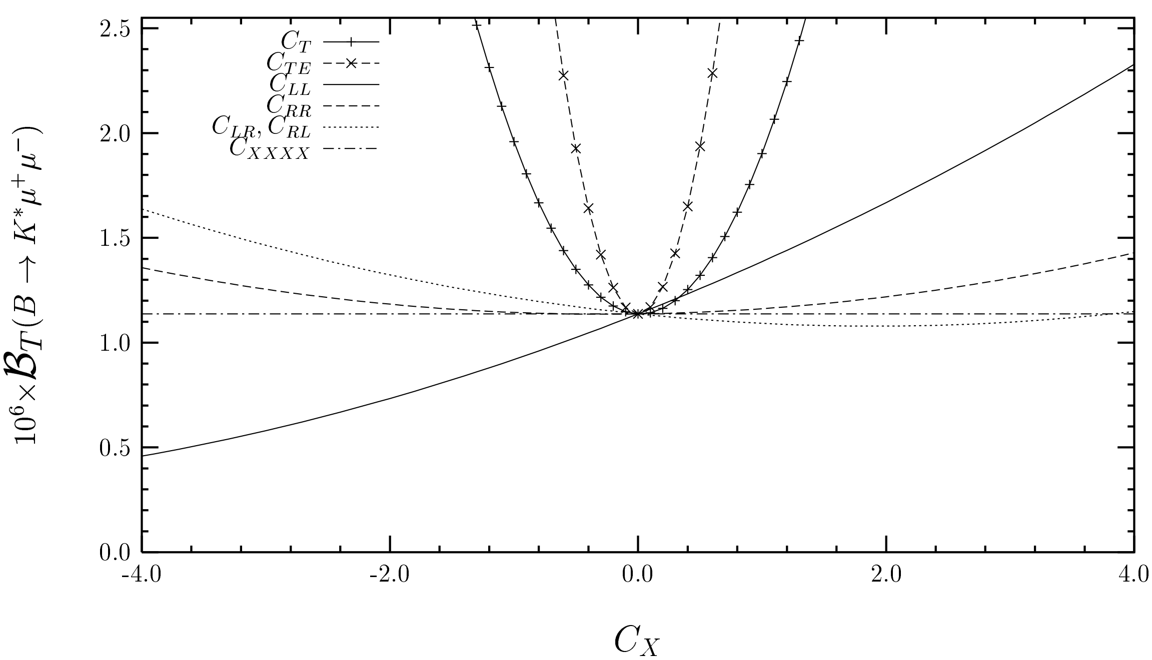

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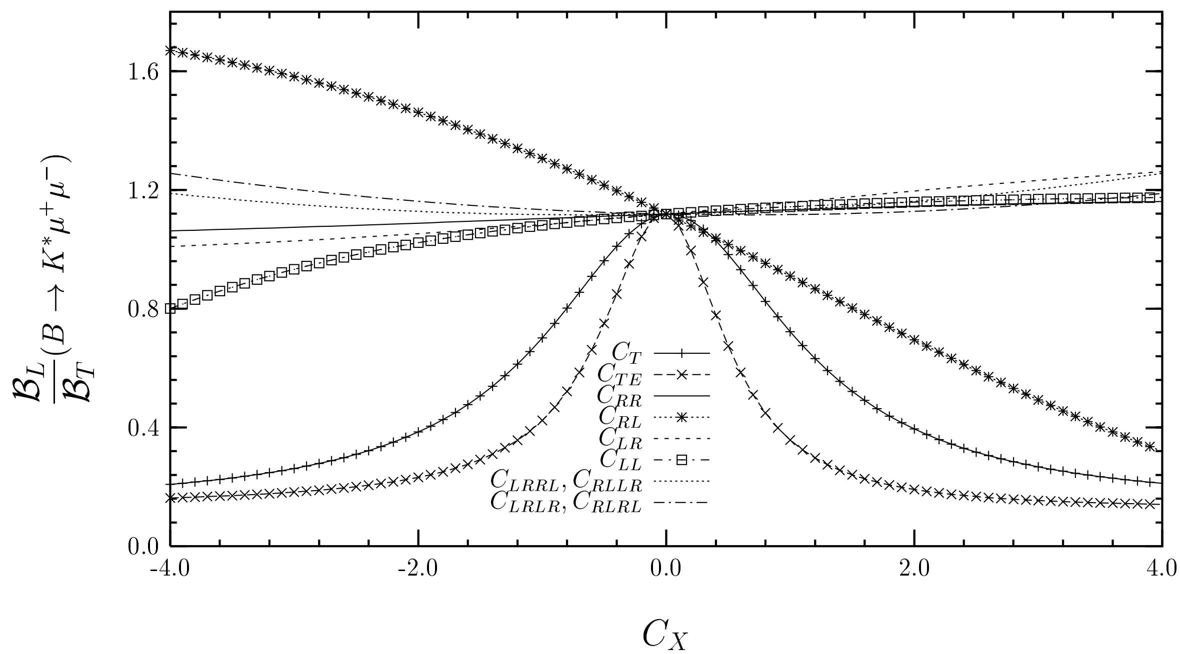

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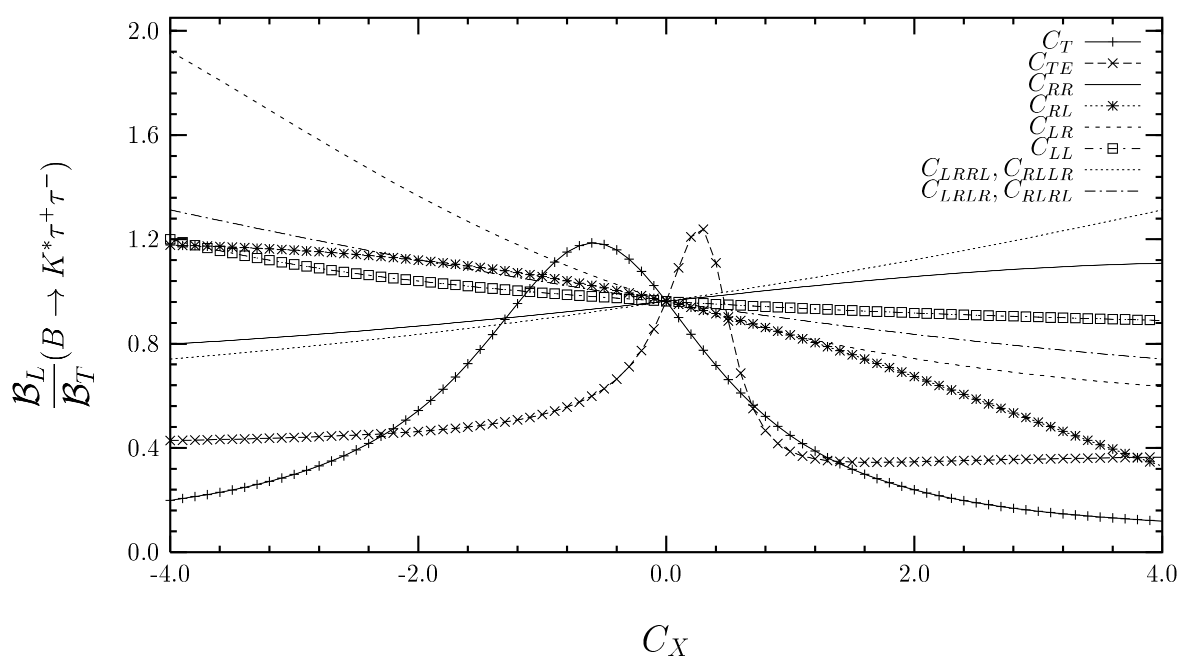

Figure 6: