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Polarized lepton pair forward-backward asymmetries in $B \to K^* \ell^+ \ell^-$ decay beyond the standard model

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Abstract: We study the polarized lepton pair forward-backward asymmetries in $B \to K^* \ell^+ \ell^-$ decay using a general, model independent form of the effective hamiltonian. We present the general expression for nine double-polarization forward-backward asymmetries. It is shown that, the study of the forward-backward asymmetries of the doubly-polarized lepton pair is a very useful tool for establishing new physics beyond the standard model.

Keywords: Rare Decays, B-Physics.

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1. Introduction

Rare B meson decays, induced by flavor changing neutral current (FCNC) $b \to s(d)\ell^+\ell^-$ transitions provide a promising ground for testing the structure of weak interactions. These decays which are forbidden in the standard model (SM) at tree level, occur at loop level and are very sensitive to the gauge structure of the SM. Moreover, these decays are also quite sensitive to the existence of new physics beyond the SM, since loops with new particles can give considerable contribution to rare decays. The new physics effects in rare decays can appear in two ways; one via modification of the existing Wilson coefficients in the SM, or through the introduction of some operators with new coefficients. Theoretical investigation of the $B \to X_s \ell^+\ell^-$ decays are relatively more clean compared to their exclusive counterparts, since they are not spoiled by nonperturbative long distance effects, while the corresponding exclusive channels are easier to measure experimentally. Some of the most important exclusive FCNC decays are $B \to K^*\gamma$ and $B \to (\pi, \rho, K, K^*)\ell^+\ell^-$ decays. The latter provides potentially a very rich set of experimental observables, such as, lepton pair forward-backward (FB) asymmetry, lepton polarizations, etc. Various kinematical distributions of such processes as $B \to K(K^*)\ell^+\ell^-$ [1, 2, 3], $B \to \pi(\rho)\ell^+\ell^-$ [4], $B_{s,d} \to \ell^+\ell^-$ [5] and $B_{s,d} \to \gamma\ell^+\ell^-$ [6] have already been studied. Experimentally measurable quantities such as forward-backward asymmetry, single polarization asymmetry, etc., have been studied for the $B \to K^*\ell^+\ell^-$ decay in [1, 7, 8, 9]. Study of these quantities can give useful information in fitting the parameters of the SM and put constraints on new physics [10, 11, 12]. It has been pointed out in [13] that the study of simultaneous polarizations of both leptons in the final state provide, in principle, measurement of many more observables which would be useful in further improvement of the parameters of the SM probing new physics beyond the SM. It should be noted here that both lepton polarizations in the $B \to K^*\tau^+\tau^-$ and $B \to K\ell^+\ell^-$ decays are studied in [14] and [15], respectively. As has already been noted, one efficient way of establishing new physics effects is studying forward-backward asymmetry in semileptonic $B \to K^*\ell^+\ell^-$ decay, since, $\mathcal{A}_{FB}$ vanishes at specific values of the dilepton invariant mass, and more essential than that, this zero position of $\mathcal{A}_{FB}$ is known to be practically free of hadronic uncertainties [12].

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The aim of the present work is studying the polarized forward-backward asymmetry in the exclusive $B \to K^* \ell^+ \ell^-$ decay using a general form of the effective hamiltonian, including all possible forms of interactions. Here we would like to remind the reader that the influence of new Wilson coefficients on various kinematical variables, such as branching ratios, lepton pair forward-backward asymmetries and single lepton polarization asymmetries for the inclusive $B \to X_{s(d)} \ell^+ \ell^-$ decays (see first references in [11, 13, 16]) and exclusive $B \to K \ell^+ \ell^-$, $K^* \ell^+ \ell^-$, $\gamma \ell^+ \ell^-$, $\pi \ell^+ \ell^-$, $\rho \ell^+ \ell^-$ [1, 2, 6, 9, 17, 18] and pure leptonic $B \to \ell^+ \ell^-$ decays [5, 19] have been studied comprehensively.

Recently, exiting results have been announced by the BaBar and Belle Collaborations for experimental study of the $B \to K^* \ell^+ \ell^-$ decay. As far as the results for the branching ratio of the $B \to K^* \ell^+ \ell^-$ decay measured by these Collaborations are gives as

$$\mathcal{B}(B \to K^* \ell^+ \ell^-) = \begin{cases} \left(11.5^{+2.6}_{-2.4} \pm 0.8 \pm 0.2\right) \times 10^{-7} & [20]\,, \\ \left(0.88^{+0.33}_{-0.29}\right) \times 10^{-6} & [21]\,. \end{cases} \tag{1.1}$$

The paper is organized as follows. In section 2, using a general form of the effective hamiltonian, we obtain the matrix element in terms of the form factors of the $B \to K^*$ transition. In section 3 we derive the analytical results for the polarized forward-backward asymmetry. Last section is devoted to the numerical analysis, discussion and conclusions.

2. Matrix element for the $B \to K^* \ell^+ \ell^-$ decay

In this section we present the matrix element for the $B \to K^* \ell^+ \ell^-$ decay using a general form of the effective hamiltonian. The $B \to K^* \ell^+ \ell^-$ process is governed by $b \to s \ell^+ \ell^-$ transition at quark level. The effective hamiltonian for the $b \to s \ell^+ \ell^-$ can be written in terms of the twelve model independent four-Fermi interactions in the following form:

$$\begin{aligned} \mathcal{H}_{\rm eff} = \frac{G_F \alpha}{\sqrt{2}\pi} V_{ts} V_{tb}^* \Bigg\{ & C_{SL}\, \bar{s} i \sigma_{\mu\nu} \frac{q^\nu}{q^2} L\, b\, \bar{\ell} \gamma^\mu \ell + C_{BR}\, \bar{s} i \sigma_{\mu\nu} \frac{q^\nu}{q^2} R\, b\, \bar{\ell} \gamma^\mu \ell + \\ & + C_{LL}^{\rm tot}\, \bar{s}_L \gamma_\mu b_L\, \bar{\ell}_L \gamma^\mu \ell_L + C_{LR}^{\rm tot}\, \bar{s}_L \gamma_\mu b_L\, \bar{\ell}_R \gamma^\mu \ell_R + C_{RL}\, \bar{s}_R \gamma_\mu b_R\, \bar{\ell}_L \gamma^\mu \ell_L + \\ & + C_{RR}\, \bar{s}_R \gamma_\mu b_R\, \bar{\ell}_R \gamma^\mu \ell_R + C_{LRLR}\, \bar{s}_L b_R\, \bar{\ell}_L \ell_R + C_{RLLR}\, \bar{s}_R b_L\, \bar{\ell}_L \ell_R + \\ & + C_{LRRL}\, \bar{s}_L b_R\, \bar{\ell}_R \ell_L + C_{RLRL}\, \bar{s}_R b_L\, \bar{\ell}_R \ell_L + C_T\, \bar{s} \sigma_{\mu\nu} b\, \bar{\ell} \sigma^{\mu\nu} \ell + \\ & + i C_{TE}\, \epsilon^{\mu\nu\alpha\beta} \bar{s} \sigma_{\mu\nu} b\, \bar{\ell} \sigma_{\alpha\beta} \ell \Bigg\}\,, \end{aligned} \tag{2.1}$$

where L and R in (2.1) 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. Here, few words about the above hamiltonian are in order. In principle, $\mathcal{O}_2$, being a member of the standard model operators, as well as operators of the type $\bar{s}_R b_L \bar{q}_L q_R$, where q represents a quark field, give contributions to the $b \to s \ell^+ \ell^-$ transition at one-loop level. The hamiltonian given in eq. (2.1) should be understood as an effective version of the most general one, where the above-mentioned contributions are absorbed into effective Wilson coefficients which depend on q^2 in general. The first two coefficients in eq. (2.1), C_{SL} and C_{BR}, are the nonlocal Fermi interactions, which correspond to $-2m_s C_7^{\rm eff}$ and $-2m_b C_7^{\rm eff}$ in the SM, respectively. The following four terms with coefficients C_{LL}, C_{LR}, C_{RL} and C_{RR} are the vector type interactions. Two of these interactions containing $C_{LL}^{\rm tot}$ and $C_{LR}^{\rm tot}$ do already exist in the SM in the form $(C_9^{\rm eff} - C_{10})$ and $(C_9^{\rm eff} + C_{10})$. Hence, representing $C_{LL}^{\rm tot}$ and $C_{LR}^{\rm tot}$ in the form

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

allows us to conclude that $C_{LL}^{\rm tot}$ and $C_{LR}^{\rm tot}$ describe the sum of the contributions from SM and the new physics. The terms with coefficients C_{LRLR}, C_{RLLR}, C_{LRRL} and C_{RLRL} describe the scalar type interactions. The remaining last two terms lead by the coefficients C_T and C_{TE}, obviously, describe the tensor type interactions.

The exclusive $B \to K^* \ell^+ \ell^-$ decay is described in terms of the matrix elements of the quark operators in eq. (2.1) over meson states, which can be parametrized in terms of the form factors. Obviously, the following matrix elements

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

are needed for the calculation of the $B \to K^* \ell^+ \ell^-$ decay. These matrix elements are defined as follows:

$$\begin{aligned} \langle K^*(p_{K^*},\varepsilon)|\bar s \gamma_\mu (1 \pm \gamma_5) b | B(p_B)\rangle &= -\epsilon_{\mu\nu\lambda\sigma}\varepsilon^{*\nu} p_{K^*}^\lambda q^\sigma \frac{2V(q^2)}{m_B+m_{K^*}} \pm i\varepsilon_\mu^* (m_B+m_{K^*}) \times \qquad (2.2) \\ &\quad \times A_1(q^2) \mp i(p_B+p_{K^*})_\mu (\varepsilon^* q) \frac{A_2(q^2)}{m_B+m_{K^*}} \mp \\ &\quad \mp i q_\mu \frac{2m_{K^*}}{q^2} (\varepsilon^* q) [A_3(q^2) - A_0(q^2)]\,, \qquad (2.3) \\ \langle K^*(p_{K^*},\varepsilon) \left| \bar s i \sigma_{\mu\nu} q^\nu (1 \pm \gamma_5) b \right| B(p_B)\rangle &= 4\epsilon_{\mu\nu\lambda\sigma}\varepsilon^{*\nu} p_{K^*}^\lambda q^\sigma T_1(q^2) \pm \\ &\quad \pm 2i \left[\varepsilon_\mu^* (m_B^2 - m_{K^*}^2) - (p_B+p_{K^*})_\mu (\varepsilon^* q)\right] T_2(q^2) \pm \\ &\quad \pm 2i (\varepsilon^* q) \left[q_\mu - (p_B+p_{K^*})_\mu \frac{q^2}{m_B^2 - m_{K^*}^2}\right] T_3(q^2)\,, \qquad (2.4) \\ \langle K^*(p_{K^*},\varepsilon) \left| \bar s \sigma_{\mu\nu} b \right| B(p_B)\rangle &= i\epsilon_{\mu\nu\lambda\sigma} \Big\{ -2T_1(q^2) \varepsilon^{*\lambda} (p_B+p_{K^*})^\sigma + \\ &\quad + \frac{2}{q^2}(m_B^2 - m_{K^*}^2)\Big[T_1(q^2) - T_2(q^2)\Big]\varepsilon^{*\lambda} q^\sigma - \\ &\quad - \frac{4}{q^2}\left[T_1(q^2) - T_2(q^2) - \frac{q^2}{m_B^2 - m_{K^*}^2} T_3(q^2)\right] \times \\ &\quad \times (\varepsilon^* q) p_{K^*}^\lambda q^\sigma \Big\}\,, \qquad (2.5) \end{aligned}$$

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 (2.3) 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 $\langle K^* | \bar{s}(1 \pm \gamma_5) b | B \rangle$ can be calculated from eq. (2.3) by contracting both sides of eq. (2.3) with q^μ and using equation of motion. Neglecting the mass of the strange quark we get

$$\langle K^*(p_{K^*}, \varepsilon) | \bar{s}(1 \pm \gamma_5) b | B(p_B) \rangle = \frac{1}{m_b} \Big[\mp 2 i m_{K^*} (\varepsilon^* q) A_0(q^2) \Big] . \tag{2.6}$$

In deriving eq. (2.6) we have used the relationship

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

which follows from the equations of motion.

Using the definition of the form factors, as given above, the amplitude of the $B \to K^* \ell^+ \ell^-$ decay can be written as

$$\begin{aligned} \mathcal{M}(B \to K^* \ell^+ \ell^-) = {} & \frac{G\alpha}{4\sqrt{2}\pi} V_{tb} V_{ts}^* \\ & \times \Bigg\{ \bar{\ell}\gamma^\mu (1-\gamma_5)\ell \Big[-2A_1 \epsilon_{\mu\nu\lambda\sigma} \varepsilon^{*\nu} p_{K^*}^\lambda q^\sigma - i B_1 \varepsilon_\mu^* + \\ & \qquad + i B_2 (\varepsilon^* q)(p_B + p_{K^*})_\mu + i B_3 (\varepsilon^* q) q_\mu \Big] + \\ & + \bar{\ell}\gamma^\mu (1+\gamma_5)\ell \Big[-2C_1 \epsilon_{\mu\nu\lambda\sigma} \varepsilon^{*\nu} p_{K^*}^\lambda q^\sigma - i D_1 \varepsilon_\mu^* + \\ & \qquad + i D_2 (\varepsilon^* q)(p_B + p_{K^*})_\mu + i D_3 (\varepsilon^* q) q_\mu \Big] \\ & + \bar{\ell}(1-\gamma_5)\ell \Big[i B_4 (\varepsilon^* q) \Big] + \bar{\ell}(1+\gamma_5)\ell \Big[i B_5 (\varepsilon^* q) \Big] + \\ & + 4 \bar{\ell} \sigma^{\mu\nu} \ell (i C_T \epsilon_{\mu\nu\lambda\sigma}) \Big[-2 T_1 \varepsilon^{*\lambda} (p_B + p_{K^*})^\sigma + \\ & \qquad + B_6 \varepsilon^{*\lambda} q^\sigma - B_7 (\varepsilon^* q) {p_{K^*}}^\lambda q^\sigma \Big] + \\ & + 16 C_{TE} \bar{\ell} \sigma_{\mu\nu} \ell \Big[-2 T_1 \varepsilon^{*\mu} (p_B + p_{K^*})^\nu + B_6 \varepsilon^{*\mu} q^\nu - \\ & \qquad - B_7 (\varepsilon^* q) {p_{K^*}}^\mu q^\nu \Big] \Bigg\} , \end{aligned} \tag{2.7}$$

where

$$\begin{aligned} A_1 &= (C_{LL}^{\text{tot}} + C_{RL}) \frac{V}{m_B + m_{K^*}} - 2(C_{BR} + C_{SL}) \frac{T_1}{q^2} , \\ B_1 &= (C_{LL}^{\text{tot}} - C_{RL})(m_B + m_{K^*}) A_1 - 2(C_{BR} - C_{SL})(m_B^2 - m_{K^*}^2) \frac{T_2}{q^2} , \\ B_2 &= \frac{C_{LL}^{\text{tot}} - C_{RL}}{m_B + m_{K^*}} A_2 - 2(C_{BR} - C_{SL}) \frac{1}{q^2} \left[T_2 + \frac{q^2}{m_B^2 - m_{K^*}^2} T_3 \right] , \\ B_3 &= 2(C_{LL}^{\text{tot}} - C_{RL}) m_{K^*} \frac{A_3 - A_0}{q^2} + 2(C_{BR} - C_{SL}) \frac{T_3}{q^2} , \\ C_1 &= A_1(C_{LL}^{\text{tot}} \to C_{LR}^{\text{tot}}, C_{RL} \to C_{RR}) , \\ D_1 &= B_1(C_{LL}^{\text{tot}} \to C_{LR}^{\text{tot}}, C_{RL} \to C_{RR}) , \\ D_2 &= B_2(C_{LL}^{\text{tot}} \to C_{LR}^{\text{tot}}, C_{RL} \to C_{RR}) , \end{aligned}$$

$$
\begin{aligned}
D_3 &= B_3(C_{LL}^{\rm tot} \to C_{LR}^{\rm tot},\, C_{RL} \to C_{RR})\,,\\
B_4 &= -2(C_{LRRL} - C_{RLRL})\frac{m_{K^*}}{m_b}A_0\,,\\
B_5 &= -2(C_{LRLR} - C_{RLLR})\frac{m_{K^*}}{m_b}A_0\,,\\
B_6 &= 2(m_B^2 - m_{K^*}^2)\frac{T_1 - T_2}{q^2}\,,\\
B_7 &= \frac{4}{q^2}\left(T_1 - T_2 - \frac{q^2}{m_B^2 - m_{K^*}^2}T_3\right).
\end{aligned}
\tag{2.8}
$$

From this expression of the decay amplitude, for the differential decay width we get the following result:

$$
\frac{d\Gamma}{d\hat{s}}(B \to K^*\ell^+\ell^-) = \frac{G^2\alpha^2 m_B}{2^{14}\pi^5}\left|V_{tb}V_{ts}^*\right|^2 \lambda^{1/2}(1,\hat{r},\hat{s})v\Delta(\hat{s})\,,
\tag{2.9}
$$

with

$$
\begin{aligned}
\Delta = \frac{2}{3\hat{r}_{K^*}\hat{s}}m_B^2\,{\rm Re}\Big[&- 6m_B\hat{m}_\ell\hat{s}\lambda(B_1 - D_1)(B_4^* - B_5^*) - \\
&- 12m_B^2\hat{m}_\ell^2\hat{s}\lambda\Big\{B_4B_5^* + (B_3 - D_2 - D_3)B_1^* - (B_2 + B_3 - D_3)D_1^*\Big\} + \\
&+ 6m_B^3\hat{m}_\ell\hat{s}(1 - \hat{r}_{K^*})\lambda(B_2 - D_2)(B_4^* - B_5^*) + \\
&+ 12m_B^4\hat{m}_\ell^2\hat{s}(1 - \hat{r}_{K^*})\lambda(B_2 - D_2)(B_3^* - D_3^*) + \\
&+ 6m_B^3\hat{m}_\ell\lambda\hat{s}^2(B_4 - B_5)(B_3^* - D_3^*) + \\
&+ 48\hat{m}_\ell^2\hat{r}_{K^*}\hat{s}\Big\{3B_1D_1^* + 2m_B^4\lambda A_1C_1^*\Big\} + \\
&+ 48m_B^5\hat{m}_\ell\hat{s}\lambda^2(B_2 + D_2)B_7^*C_{TE}^* - \\
&- 16m_B^4\hat{r}_{K^*}\hat{s}(\hat{m}_\ell^2 - \hat{s})\lambda\Big\{\left|A_1\right|^2 + \left|C_1\right|^2\Big\} - \\
&- m_B^2\hat{s}(2\hat{m}_\ell^2 - \hat{s})\lambda\Big\{\left|B_4\right|^2 + \left|B_5\right|^2\Big\} - \\
&- 48m_B^3\hat{m}_\ell\hat{s}(1 - \hat{r}_{K^*} - \hat{s})\lambda\Big\{(B_1 + D_1)B_7^*C_{TE}^* + 2(B_2 + D_2)B_6^*C_{TE}^*\Big\} - \\
&- 6m_B^4\hat{m}_\ell^2\hat{s}\lambda\Big\{2(2 + 2\hat{r}_{K^*} - \hat{s})B_2D_2^* - \hat{s}\left|(B_3 - D_3)\right|^2\Big\} + \\
&+ 96m_B\hat{m}_\ell\hat{s}(\lambda + 12\hat{r}_{K^*}\hat{s})(B_1 + D_1)B_6^*C_{TE}^* + \\
&+ 8m_B^2\hat{s}^2\Big\{v^2\left|C_T\right|^2 + 4(3 - 2v^2)\left|C_{TE}\right|^2\Big\} \times \\
&\times \Big\{4(\lambda + 12\hat{r}_{K^*}\hat{s})\left|B_6\right|^2 - 4m_B^2\lambda(1 - \hat{r}_{K^*} - \hat{s})B_6B_7^* + m_B^4\lambda^2\left|B_7\right|^2\Big\} - \\
&- 4m_B^2\lambda\Big\{\hat{m}_\ell^2(2 - 2\hat{r}_{K^*} + \hat{s}) + \hat{s}(1 - \hat{r}_{K^*} - \hat{s})\Big\}(B_1B_2^* + D_1D_2^*) + \\
&+ \hat{s}\Big\{6\hat{r}_{K^*}\hat{s}(3 + v^2) + \lambda(3 - v^2)\Big\}\Big\{\left|B_1\right|^2 + \left|D_1\right|^2\Big\} - \\
&- 2m_B^4\lambda\Big\{\hat{m}_\ell^2[\lambda - 3(1 - \hat{r}_{K^*})^2] - \lambda\hat{s}\Big\}\Big\{\left|B_2\right|^2 + \left|D_2\right|^2\Big\} + \\
&+ 128m_B^2\Big\{4\hat{m}_\ell^2[20\hat{r}_{K^*}\lambda - 12\hat{r}_{K^*}(1 - \hat{r}_{K^*})^2 - \lambda\hat{s}] + \\
&\qquad + \hat{s}[4\hat{r}_{K^*}\lambda + 12\hat{r}_{K^*}(1 - \hat{r}_{K^*})^2 + \lambda\hat{s}]\Big\}\left|C_T\right|^2\left|T_1\right|^2 +
\end{aligned}
$$

$$
\begin{aligned}
&+512m_B^2\Big\{\hat{s}[4\hat{r}_{K^*}\lambda+12\hat{r}_{K^*}(1-\hat{r}_{K^*})^2+\lambda\hat{s}]+\\
&+8\hat{m}_\ell^2[12\hat{r}_{K^*}(1-\hat{r}_{K^*})^2+\lambda(\hat{s}-8\hat{r}_{K^*})]\Big\}\left|C_{TE}\right|^2\left|T_1\right|^2-\\
&-64m_B^2\hat{s}^2\Big\{v^2\left|C_T\right|^2+4(3-2v^2)\left|C_{TE}\right|^2\Big\}\times\\
&\times\Big\{2[\lambda+12\hat{r}_{K^*}(1-\hat{r}_{K^*})]B_6T_1^*-m_B^2\lambda(1+3\hat{r}_{K^*}-\hat{s})B_7T_1^*\Big\}+\\
&+768m_B^3\hat{m}_\ell\hat{r}_{K^*}\hat{s}\lambda(A_1+C_1)C_T^*T_1^*-\\
&-192m_B\hat{m}_\ell\hat{s}[\lambda+12\hat{r}_{K^*}(1-\hat{r}_{K^*})](B_1+D_1)C_{TE}^*T_1^*+\\
&+192m_B^3\hat{m}_\ell\hat{s}\lambda(1+3\hat{r}_{K^*}-\hat{s})\lambda(B_2+D_2)C_{TE}^*T_1^*\Big]\,, \qquad (2.10)
\end{aligned}
$$

where $\hat{s}=q^2/m_B^2$, $\hat{r}=m_{K^*}^2/m_B^2$ and $\lambda(a,b,c)=a^2+b^2+c^2-2ab-2ac-2bc$, $\hat{m}_\ell=m_\ell/m_B$, $v=\sqrt{1-4\hat{m}_\ell^2/\hat{s}}$ is the final lepton velocity.

The definition of the polarized FB asymmetries will be presented in the next section.

3. Polarized forward-backward asymmetries of leptons

In this section we calculate the polarized FB asymmetries. For this purpose, we define the following orthogonal unit vectors $s_i^{\pm\mu}$ in the rest frame of $\ell^\pm$, where $i=L,N$ or T correspond to longitudinal, normal, transversal polarization directions, respectively (see also [1, 8, 10, 14]),

$$
\begin{aligned}
s_L^{-\mu} &= \left(0,\vec{e}_L^{\,-}\right)=\left(0,\frac{\vec{p}_-}{|\vec{p}_-|}\right)\,,\\
s_N^{-\mu} &= \left(0,\vec{e}_N^{\,-}\right)=\left(0,\frac{\vec{p}_K\times\vec{p}_-}{|\vec{p}_K\times\vec{p}_-|}\right)\,,\\
s_T^{-\mu} &= \left(0,\vec{e}_T^{\,-}\right)=\left(0,\vec{e}_N^{\,-}\times\vec{e}_L^{\,-}\right)\,,\\
s_L^{+\mu} &= \left(0,\vec{e}_L^{\,+}\right)=\left(0,\frac{\vec{p}_+}{|\vec{p}_+|}\right)\,,\\
s_N^{+\mu} &= \left(0,\vec{e}_N^{\,+}\right)=\left(0,\frac{\vec{p}_K\times\vec{p}_+}{|\vec{p}_K\times\vec{p}_+|}\right)\,,\\
s_T^{+\mu} &= \left(0,\vec{e}_T^{\,+}\right)=\left(0,\vec{e}_N^{\,+}\times\vec{e}_L^{\,+}\right)\,, \qquad (3.1)
\end{aligned}
$$

where $\vec{p}_\mp$ and $\vec{p}_K$ are the three-momenta of the leptons $\ell^\mp$ and K^* meson in the center of mass frame (CM) of $\ell^-\,\ell^+$ system, respectively. Transformation of unit vectors from the rest frame of the leptons to CM frame of leptons can be accomplished by the Lorentz boost. Boosting of the longitudinal unit vectors $s_L^{\pm\mu}$ yields

$$
\left(s_L^{\mp\mu}\right)_{CM}=\left(\frac{|\vec{p}_\mp|}{m_\ell},\frac{E_\ell\vec{p}_\mp}{m_\ell\,|\vec{p}_\mp|}\right)\,, \qquad (3.2)
$$

where $\vec{p}_+=-\vec{p}_-$, E_ℓ and m_ℓ are the energy and mass of leptons in the CM frame, respectively. The remaining two unit vectors $s_N^{\pm\mu}$, $s_T^{\pm\mu}$ are unchanged under Lorentz boost.

The definition of the unpolarized and normalized differential forward-backward asymmetry is (see for example [22])

$$\mathcal{A}_{FB} = \frac{\int_0^1 \frac{d^2\Gamma}{d\hat{s}dz} - \int_{-1}^0 \frac{d^2\Gamma}{d\hat{s}dz}}{\int_0^1 \frac{d^2\Gamma}{d\hat{s}dz} + \int_{-1}^0 \frac{d^2\Gamma}{d\hat{s}dz}} \,, \tag{3.3}$$

where $z = \cos\theta$ is the angle between B meson and ℓ^- in the center mass frame of leptons. When the spins of both leptons are taken into account, the $\mathcal{A}_{FB}$ will be a function of the spins of the final leptons and it is defined as

$$\begin{aligned}
\mathcal{A}_{FB}^{ij}(\hat{s}) &= \left(\frac{d\Gamma(\hat{s})}{d\hat{s}}\right)^{-1} \left\{ \int_0^1 dz - \int_{-1}^0 dz \right\} \times \\
&\quad \times \left\{ \left[\frac{d^2\Gamma(\hat{s}, \vec{s}^- = \vec{i}, \vec{s}^+ = \vec{j})}{d\hat{s}dz} - \frac{d^2\Gamma(\hat{s}, \vec{s}^- = \vec{i}, \vec{s}^+ = -\vec{j})}{d\hat{s}dz} \right] - \right. \\
&\quad \left. - \left[\frac{d^2\Gamma(\hat{s}, \vec{s}^- = -\vec{i}, \vec{s}^+ = \vec{j})}{d\hat{s}dz} - \frac{d^2\Gamma(\hat{s}, \vec{s}^- = -\vec{i}, \vec{s}^+ = -\vec{j})}{d\hat{s}dz} \right] \right\}, \\
&= \mathcal{A}_{FB}(\vec{s}^- = \vec{i}, \vec{s}^+ = \vec{j}) - \mathcal{A}_{FB}(\vec{s}^- = \vec{i}, \vec{s}^+ = -\vec{j}) - \mathcal{A}_{FB}(\vec{s}^- = -\vec{i}, \vec{s}^+ = \vec{j}) + \\
&\quad + \mathcal{A}_{FB}(\vec{s}^- = -\vec{i}, \vec{s}^+ = -\vec{j}) \,.
\end{aligned} \tag{3.4}$$

Using these definitions for the double polarized FB asymmetries, we get the following results:

$$\begin{aligned}
\mathcal{A}_{FB}^{LL} &= \frac{2}{\hat{r}_{K^*}\Delta} m_B^3 \sqrt{\lambda} v \,\mathrm{Re} \Big[- m_B^3 \hat{m}_\ell \lambda \Big\{ 4(B_1 - D_1) B_7^* C_T^* - (B_4 + B_5)(B_2^* + D_2^*) \Big\} + \\
&\quad + 4 m_B^4 \hat{m}_\ell \lambda \Big\{ (1 - \hat{r}_{K^*})(B_2 - D_2) B_7^* C_T^* + \hat{s}(B_3 - D_3) B_7^* C_T^* \Big\} - \\
&\quad - \hat{m}_\ell (1 - \hat{r}_{K^*} - \hat{s}) \Big\{ B_1^* (B_4 + B_5 - 8 B_6 C_T) + \\
&\qquad + D_1^* (B_4 + B_5 + 8 B_6 C_T) \Big\} + \\
&\quad + 8 m_B \hat{r}_{K^*} \hat{s} (A_1 B_1^* - C_1 D_1^*) + \\
&\quad + 128 m_B^2 \hat{m}_\ell \hat{r}_{K^*} \hat{s} (A_1 - C_1) B_6^* C_{TE}^* + \\
&\quad + 2 m_B^3 \hat{s} \lambda \Big\{ (B_4 - B_5) B_7^* C_T^* + 2(B_4 + B_5) B_7^* C_{TE}^* \Big\} - \\
&\quad - 8 m_B^2 \hat{m}_\ell (1 - \hat{r}_{K^*})(1 - \hat{r}_{K^*} - \hat{s})(B_2 - D_2) B_6^* C_T^* - \\
&\quad - 4 m_B (1 - \hat{r}_{K^*} - \hat{s}) \hat{s} \Big\{ (B_4 - B_5) B_6^* C_T^* + 2(B_4 + B_5) B_6^* C_{TE}^* + \\
&\qquad + 2 m_B \hat{m}_\ell (B_3 - D_3) B_6^* C_T^* \Big\} - \\
&\quad - 256 m_B^5 \hat{m}_\ell \hat{r}_{K^*} (1 - \hat{r}_{K^*})(A_1 - C_1) T_1^* C_{TE}^* - \\
&\quad - 16 \hat{m}_\ell (1 - 5\hat{r}_{K^*} - \hat{s})(B_1 - D_1) T_1^* C_T^* + \\
&\quad + 16 m_B^2 \hat{m}_\ell (1 - \hat{r}_{K^*})(1 + 3\hat{r}_{K^*} - \hat{s})(B_2 - D_2) T_1^* C_T^* + \\
&\quad + 8 m_B (1 + 3\hat{r}_{K^*} - \hat{s}) \hat{s} \Big\{ 2(B_4 + B_5) T_1^* C_{TE}^* + (B_4 - B_5) T_1^* C_T^* + \\
&\qquad + 2 m_B \hat{m}_\ell (B_3 - D_3) T_1^* C_T^* \Big\} \Big] \,,
\end{aligned} \tag{3.5}$$

$$\mathcal{A}_{FB}^{LN} = \frac{8}{3\hat{r}_{K^*}\hat{s}\Delta} m_B^2 \sqrt{\hat{s}} \lambda v \,\mathrm{Im} \Big[- \hat{m}_\ell (B_1 D_1^* + m_B^4 \lambda B_2 D_2^*) + 4 m_B^4 \hat{m}_\ell \hat{r}_{K^*} \sqrt{\hat{s}} A_1 C_1^* -$$

$$
\begin{aligned}
&- 2m_B\hat{s}\Big\{B_6(C_T - 2C_{TE})B_1^* + B_6(C_T + 2C_{TE})D_1^*\Big\} - \\
&- m_B^5\hat{s}\lambda\Big\{B_7(C_T - 2C_{TE})B_2^* + B_7(C_T + 2C_{TE})D_2^*\Big\} - \\
&- 16m_B^2\hat{m}_\ell\hat{s}\Big(4\left|B_6\right|^2 + m_B^4\lambda\left|B_7\right|^2\Big)C_T C_{TE}^* + \\
&+ m_B^2\hat{m}_\ell(1-\hat{r}_{K^*}-\hat{s})(B_1D_2^* + B_2D_1^*) + \\
&+ m_B^3\hat{s}(1-\hat{r}_{K^*}-\hat{s})\Big\{(B_1^*B_7 + 2B_2^*B_6)(C_T - 2C_{TE}) + \\
&\qquad\qquad + (D_1^*B_7 + 2D_2^*B_6)(C_T + 2C_{TE})\Big\} - \\
&- 64m_B^2\hat{m}_\ell\hat{s}\Big\{ - m_B^2(1-\hat{r}_{K^*}-\hat{s})\,\mathrm{Re}[B_6B_7^*] + 4\left|T_1\right|^2 - \\
&\qquad - 4\,\mathrm{Re}[B_6T_1^*] + 2m_B^2(1+3\hat{r}_{K^*}-\hat{s})\,\mathrm{Re}[B_7T_1^*]\Big\} \times \\
&\times C_T C_{TE}^* + 16m_B^3\hat{r}_{K^*}\hat{s}\Big\{(A_1 - C_1)C_T^*T_1^* - 2(A_1 + C_1)C_{TE}^*T_1^*\Big\} + \\
&+ 4m_B\hat{s}\Big\{B_1^*(C_T - 2C_{TE})T_1 + D_1^*(C_T + 2C_{TE})T_1\Big\} - \\
&- 4m_B^3\hat{s}(1+3\hat{r}_{K^*}-\hat{s})\Big\{B_2^*(C_T - 2C_{TE})T_1 + D_2^*(C_T + 2C_{TE})T_1\Big\}\Big]\,,
\end{aligned}
\tag{3.6}
$$

$$
\begin{aligned}
\mathcal{A}_{FB}^{NL} = \frac{8}{3\hat{r}_{K^*}\hat{s}\Delta}m_B^2\sqrt{\hat{s}}\lambda v\,\mathrm{Im}\Big[&- \hat{m}_\ell(B_1D_1^*] + m_B^4\lambda B_2D_2^*) + 4m_B^2\hat{m}_\ell\hat{r}_{K^*}\hat{s}A_1C_1^* + \\
&+ 2m_B\hat{s}\Big\{B_6(C_T + 2C_{TE})B_1^* + B_6(C_T - 2C_{TE})D_1^*\Big\} + \\
&+ m_B^5\hat{s}\lambda\Big\{B_7(C_T + 2C_{TE})B_2^* + B_7(C_T - 2C_{TE})D_2^*\Big\} + \\
&+ 16m_B^2\hat{m}_\ell\hat{s}(4\left|B_6\right|^2 + m_B^4\lambda\left|B_7\right|^2)C_T C_{TE}^* + \\
&+ m_B^2\hat{m}_\ell(1-\hat{r}_{K^*}-\hat{s})(B_1D_2^* + B_2D_1^*) - \\
&- m_B^3\hat{s}(1-\hat{r}_{K^*}-\hat{s})\Big\{(B_1^*B_7 + 2B_2^*B_6)(C_T + 2C_{TE}) + \\
&\qquad\qquad + (D_1^*B_7 + 2D_2^*B_6)(C_T - 2C_{TE})\Big\} + \\
&+ 64m_B^2\hat{m}_\ell\hat{s}\Big\{ - m_B^2(1-\hat{r}_{K^*}-\hat{s})\,\mathrm{Re}[B_6B_7^*] + 4\left|T_1\right|^2 - 4\,\mathrm{Re}[B_6T_1^*] + \\
&\qquad + 2m_B^2(1+3\hat{r}_{K^*}-\hat{s})\,\mathrm{Re}[B_7T_1^*]\Big\}C_T C_{TE}^* + \\
&+ 16m_B^3\hat{r}_{K^*}\hat{s}\Big\{(A_1 - C_1)C_T^*T_1^* + 2(A_1 + C_1)C_{TE}^*T_1^*\Big\} - \\
&- 4m_B\hat{s}\Big\{B_1^*(C_T + 2C_{TE})T_1 + D_1^*(C_T - 2C_{TE})T_1\Big\} + \\
&+ 4m_B^3\hat{s}(1+3\hat{r}_{K^*}-\hat{s})\Big\{B_2^*(C_T + 2C_{TE})T_1 + \\
&\qquad\qquad + D_2^*(C_T - 2C_{TE})T_1\Big\}\Big]\,,
\end{aligned}
\tag{3.7}
$$

$$
\begin{aligned}
\mathcal{A}_{FB}^{LT} = \frac{4}{3\hat{r}_{K^*}\hat{s}\Delta}m_B^2\sqrt{\hat{s}}\lambda\,\mathrm{Re}\Big[&- \hat{m}_\ell\Big\{\left|B_1 + D_1\right|^2 + m_B^4\lambda\left|B_2 + D_2\right|^2\Big\} \\
&+ 4m_B^4\hat{m}_\ell\hat{r}_{K^*}\hat{s}\Big\{\left|A_1 + C_1\right|^2\Big\} - \\
&- 64m_B^2\hat{m}_\ell\hat{s}\left|C_{TE}\right|^2\Big\{4\left|B_6\right|^2 + m_B^4\lambda\left|B_7\right|^2 - \\
&\qquad\qquad - 4m_B^2(1-\hat{r}_{K^*}-\hat{s})B_6B_7^*\Big\} + \\
&+ 2m_B^2\hat{m}_\ell(1-\hat{r}_{K^*}-\hat{s})(B_1 + D_1)(B_2^* + D_2^*) + \\
&+ 2m_B^3(1-\hat{r}_{K^*}-\hat{s})\Big\{4\hat{m}_\ell^2(2B_2^*B_6 + B_1^*B_7)(C_T + 2C_{TE}) -
\end{aligned}
$$

$$
\begin{aligned}
& - \hat{s}(2B_2^* B_6 + B_1^* B_7)(C_T - 2C_{TE})\Big\} - \\
& - 4m_B\Big\{4\hat{m}_\ell^2\Big[B_1^* B_6(C_T + 2C_{TE}) - B_6 D_1^*(C_T - 2C_{TE})\Big] - \\
& - \hat{s}\Big[B_1^* B_6(C_T - 2C_{TE}) - B_6 D_1^*(C_T + 2C_{TE})\Big]\Big\} - \\
& - 2m_B^5\lambda\Big\{4\hat{m}_\ell^2\Big[B_2^* B_7(C_T + 2C_{TE}) - B_7 D_2^*(C_T - 2C_{TE})\Big] - \\
& - \hat{s}\Big[B_2^* B_7(C_T - 2C_{TE}) - B_7 D_2^*(C_T + 2C_{TE})\Big]\Big\} - \\
& - 2m_B^3(1 - \hat{r}_{K^*} - \hat{s})\Big\{4\hat{m}_\ell^2(2B_6 D_2^* + B_7 D_1^*)(C_T - 2C_{TE}) - \\
& - \hat{s}(2B_6 D_2^* + B_7 D_1^*)(C_T + 2C_{TE})\Big\} + \\
& + 256 m_B^2 \hat{m}_\ell\Big\{2\hat{s}\left|C_{TE}\right|^2\Big[2B_6 T_1^* - m_B^2(1 + 3\hat{r}_{K^*} - \hat{s})B_7 T_1^*\Big] + \\
& + 4\left|T_1\right|^2\Big[\hat{r}_{K^*}\left|C_T\right|^2 + (4\hat{r}_{K^*} - \hat{s})\left|C_{TE}\right|^2\Big]\Big\} + \\
& + 32 m_B^3 \hat{r}_{K^*}\Big\{4\hat{m}_\ell^2\Big[(A_1 + C_1)C_T^* T_1^* + 2(A_1 - C_1)C_{TE}^* T_1^*\Big] + \\
& + \hat{s}\Big[A_1^*(C_T - 2C_{TE})T_1 + C_1^*(C_T + 2C_{TE})T_1\Big]\Big\} + \\
& + 8m_B\Big\{4\hat{m}_\ell^2(C_T + 2C_{TE}) - \hat{s}(C_T - 2C_{TE})\Big\} \times \\
& \times \Big\{B_1^* - m_B^2(1 + 3\hat{r}_{K^*} - \hat{s})B_2^*\Big\} T_1 \\
& - 8m_B\Big\{4\hat{m}_\ell^2(C_T - 2C_{TE}) - \hat{s}(C_T + 2C_{TE})\Big\} \times \\
& \times \Big\{D_1^* - m_B^2(1 + 3\hat{r}_{K^*} - \hat{s})D_2^*\Big\} T_1\Big], \qquad (3.8)
\end{aligned}
$$

$$
\begin{aligned}
\mathcal{A}_{FB}^{TL} = \frac{4}{3\hat{r}_{K^*}\hat{s}\Delta} m_B^2 \sqrt{\hat{s}}\lambda\, \mathrm{Re}\Big[& \hat{m}_\ell\Big\{\left|B_1 + D_1\right|^2 + m_B^4\lambda\left|B_2 + D_2\right|^2\Big\} - \\
& - 4m_B^4 \hat{m}_\ell \hat{r}_{K^*}\Big\{\left|A_1 + C_1\right|^2\Big\} + \\
& + 64 m_B^2 \hat{m}_\ell \hat{s}\left|C_{TE}\right|^2\Big\{4\left|B_6\right|^2 + m_B^4\lambda\left|B_7\right|^2 - \\
& - 4m_B^2(1 - \hat{r}_{K^*} - \hat{s})B_6 B_7^*\Big\} - \\
& - 2m_B^2 \hat{m}_\ell(1 - \hat{r}_{K^*} - \hat{s})(B_1 + D_1)(B_2^* + D_2^*) + \\
& + 2m_B^3(1 - \hat{r}_{K^*} - \hat{s})\Big\{4\hat{m}_\ell^2(2B_2^* B_6 + B_1^* B_7)(C_T - 2C_{TE}) - \\
& - \hat{s}(2B_2^* B_6 + B_1^* B_7)(C_T + 2C_{TE})\Big\} - \\
& - 4m_B\Big\{4\hat{m}_\ell^2\Big[B_1^* B_6(C_T - 2C_{TE}) - B_6 D_1^*(C_T + 2C_{TE})\Big] - \\
& - \hat{s}\Big[B_1^* B_6(C_T + 2C_{TE}) - B_6 D_1^*(C_T - 2C_{TE})\Big]\Big\} - \\
& - 2m_B^5\lambda\Big\{4\hat{m}_\ell^2\Big[B_2^* B_7(C_T - 2C_{TE}) - B_7 D_2^*(C_T + 2C_{TE})\Big] - \\
& - \hat{s}\Big[B_2^* B_7(C_T + 2C_{TE}) - B_7 D_2^*(C_T - 2C_{TE})\Big]\Big\} - \\
& - 2m_B^3(1 - \hat{r}_{K^*} - \hat{s})\Big\{4\hat{m}_\ell^2(2B_6 D_2^* + B_7 D_1^*)(C_T + 2C_{TE}) - \\
& - \hat{s}(2B_6 D_2^* + B_7 D_1^*)(C_T - 2C_{TE})\Big\} - \\
& - 256 m_B^2 \hat{m}_\ell\Big\{2\hat{s}\left|C_{TE}\right|^2\Big[2B_6 T_1^* - m_B^2(1 + 3\hat{r}_{K^*} - \hat{s})B_7 T_1^*\Big] +
\end{aligned}
$$

$$
\begin{aligned}
&+ 4\left|T_1\right|^2\left[\hat r_{K^*}\left|C_T\right|^2 + (4\hat r_{K^*} - \hat s)\left|C_{TE}\right|^2\right]\Big\} - \\
&- 32 m_B^3 \hat r_{K^*}\Big\{4\hat m_\ell^2\Big[(A_1 + C_1)C_T^* T_1^* - 2(A_1 - C_1)C_{TE}^* T_1^*\Big] + \\
&+ \hat s\Big[A_1^*(C_T + 2C_{TE})T_1 + C_1^*(C_T - 2C_{TE})T_1\Big]\Big\} - \\
&- 8 m_B\Big\{4\hat m_\ell^2(C_T + 2C_{TE}) - \hat s(C_T - 2C_{TE})\Big\} \times \\
&\times \Big\{D_1^* - m_B^2(1 + 3\hat r_{K^*} - \hat s)D_2^*\Big\}T_1 \\
&+ 8 m_B\Big\{4\hat m_\ell^2(C_T - 2C_{TE}) - \hat s(C_T + 2C_{TE})\Big\} \times \\
&\times \Big\{B_1^* - m_B^2(1 + 3\hat r_{K^*} - \hat s)B_2^*\Big\}T_1\Big], \qquad (3.9)
\end{aligned}
$$

$$
\begin{aligned}
\mathcal{A}_{FB}^{NT} = \frac{2}{\hat r_{K^*}\hat s\Delta} m_B^2\sqrt{\lambda}\,\mathrm{Im}\Big[& m_B^3\hat m_\ell\hat s\lambda\Big\{(B_4 - B_5)(B_2^* + D_2^*) + 8B_7 C_{TE}(B_1^* - D_1^*) + \\
&+ 8m_B^2\hat s B_7^* C_{TE}^*(B_3 - D_3)\Big\} - \\
&- 2m_B^4\hat m_\ell^2\hat s\lambda(B_2 + D_2)(B_3^* - D_3^*) + \\
&+ 4m_B^4\hat m_\ell(1 - \hat r_{K^*})\lambda\Big\{2m_B\hat s B_7^* C_{TE}^*(B_2 - D_2) + \hat m_\ell B_2 D_2^*\Big\} + \\
&+ 2m_B^2\hat m_\ell^2\hat s(1 + 3\hat r_{K^*} - \hat s)(B_1 B_2^* - D_1 D_2^*) + \\
&+ \hat m_\ell(1 - \hat r_{K^*} - \hat s)\Big\{m_B\hat s\Big[- B_1^*(B_4 - B_5 + 16B_6 C_{TE}) - \\
&- D_1^*(B_4 - B_5 - 16B_6 C_{TE}) + \\
&+ 2m_B\hat m_\ell(B_1 + D_1)(B_3^* - D_3^*)\Big] + \\
&+ 4\Big[\hat m_\ell B_1 D_1^* + 4m_B^3\hat s^2 B_6 C_{TE}(B_3^* - D_3^*)\Big]\Big\} - \\
&- 16m_B^3\hat m_\ell\hat s(1 - \hat r_{K^*})(1 - \hat r_{K^*} - \hat s)(B_2 - D_2)B_6^* C_{TE}^* + \\
&+ 2m_B^2\hat m_\ell^2[\lambda + (1 - \hat r_{K^*})(1 - \hat r_{K^*} - \hat s)](B_1^* D_2 + B_2^* D_1) + \\
&+ 32m_B^3\hat m_\ell\hat s(1 - \hat r_{K^*})(1 + 3\hat r_{K^*} - \hat s)(B_2 - D_2)C_{TE}^* T_1^* - \\
&- 8m_B\hat s(1 + 3\hat r_{K^*} - \hat s)\Big\{4\hat m_\ell(B_1 - D_1)C_{TE}^* T_1^* - \\
&- 2m_B\hat s(B_4 - B_5)C_{TE}^* T_1^* - \\
&- 4m_B^2\hat m_\ell\hat s(B_3 - D_3)C_{TE}^* T_1^* + \\
&+ m_B\hat s v^2(B_4 + B_5)C_T^* T_1^*\Big\} - \\
&- 4m_B^2\hat s^2(1 - \hat r_{K^*} - \hat s)\Big\{2(B_4 - B_5)B_6^* C_{TE}^* - v^2(B_4 + B_5)B_6^* C_T^*\Big\} + \\
&+ 2m_B^4\hat s^2\lambda\Big\{2(B_4 - B_5)B_7^* C_{TE}^* - v^2(B_4 + B_5)B_7^* C_T^*\Big\}\Big], \qquad (3.10)
\end{aligned}
$$

$$
\begin{aligned}
\mathcal{A}_{FB}^{TN} = \frac{2}{\hat r_{K^*}\hat s\Delta} m_B^2\sqrt{\lambda}\,\mathrm{Im}\Big[& m_B^3\hat m_\ell\hat s\lambda\Big\{(B_4 - B_5)(B_2^* + D_2^*) + 8B_7 C_{TE}(B_1^* - D_1^*) + \\
&+ 8m_B^2\hat s B_7^* C_{TE}^*(B_3 - D_3)\Big\} - \\
&- 2m_B^4\hat m_\ell^2\hat s\lambda(B_2 + D_2)(B_3^* - D_3^*) + \\
&+ 4m_B^4\hat m_\ell(1 - \hat r_{K^*})\lambda\Big\{2m_B\hat s B_7^* C_{TE}^*(B_2 - D_2) + \hat m_\ell B_2 D_2^*\Big\} + \\
&+ 2m_B^2\hat m_\ell^2\hat s(1 + 3\hat r_{K^*} - \hat s)\,\mathrm{Im}(B_1 B_2^* - D_1 D_2^*) + \\
&+ \hat m_\ell(1 - \hat r_{K^*} - \hat s)\Big\{m_B\hat s\Big[B_1^*(B_4 - B_5 + 16B_6 C_{TE}) + \\
&+ D_1^*(B_4 - B_5 - 16B_6 C_{TE}) -
\end{aligned}
$$

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$$
\begin{aligned}
& \qquad\qquad\qquad\qquad\qquad - 2m_B\hat{m}_\ell(B_1+D_1)(B_3^*-D_3^*)\Big] + \\
& \qquad\qquad\qquad + 4\Big[\hat{m}_\ell B_1 D_1^* + 4m_B^3\hat{s}^2 B_6 C_{TE}(B_3^*-D_3^*)\Big]\Big\} - \\
& - 16m_B^3\hat{m}_\ell\hat{s}(1-\hat{r}_{K^*})(1-\hat{r}_{K^*}-\hat{s})(B_2-D_2)B_6^*C_{TE}^* + \\
& + 2m_B^2\hat{m}_\ell^2[\lambda+(1-\hat{r}_{K^*})(1-\hat{r}_{K^*}-\hat{s})](B_1^*D_2+B_2^*D_1) + \\
& + 32m_B^3\hat{m}_\ell\hat{s}(1-\hat{r}_{K^*})(1+3\hat{r}_{K^*}-\hat{s})(B_2-D_2)C_{TE}^*T_1^* \\
& - 8m_B\hat{s}(1+3\hat{r}_{K^*}-\hat{s})\Big\{4\hat{m}_\ell(B_1-D_1)C_{TE}^*T_1^* - \\
& \qquad - 2m_B\hat{s}(B_4-B_5)C_{TE}^*T_1^* - \\
& \qquad - 4m_B^2\hat{m}_\ell\hat{s}(B_3-D_3)C_{TE}^*T_1^* + \\
& \qquad + m_B\hat{s}v^2(B_4+B_5)C_T^*T_1^*\Big\} - \\
& - 4m_B^2\hat{s}^2(1-\hat{r}_{K^*}-\hat{s})\Big\{2(B_4-B_5)B_6^*C_{TE}^* - v^2(B_4+B_5)B_6^*C_T^*\Big\} + \\
& + 2m_B^4\hat{s}^2\lambda\Big\{2(B_4-B_5)B_7^*C_{TE}^* - v^2(B_4+B_5)B_7^*C_T^*\Big\}\Big],
\end{aligned} \tag{3.11}
$$

$$
\begin{aligned}
\mathcal{A}_{FB}^{NN} = \frac{2}{\hat{r}_{K^*}\Delta}m_B^3\sqrt{\lambda}v\,\mathrm{Re}\Big[& - m_B^2\hat{m}_\ell\lambda\Big\{4(B_1-D_1)B_7^*C_T^* + (B_2+D_2)(B_4^*+B_5^*)\Big\} + \\
& + 4m_B^4\hat{m}_\ell\lambda\Big\{(1-\hat{r}_{K^*})(B_2-D_2)B_7^*C_T^* + \hat{s}(B_3-D_3)B_7^*C_T^*\Big\} + \\
& + 2m_B^3\hat{s}\lambda\Big\{(B_4-B_5)B_7^*C_T^* - 2(B_4+B_5)B_7^*C_{TE}^*\Big\} + \\
& + \hat{m}_\ell(1-\hat{r}_{K^*}-\hat{s})\Big\{B_1^*(B_4+B_5+8B_6C_T) + \\
& + D_1^*(B_4+B_5-8B_6C_T)\Big\} - \\
& - 8m_B^2\hat{m}_\ell(1-\hat{r}_{K^*})(1-\hat{r}_{K^*}-\hat{s})(B_2-D_2)B_6^*C_T^* - \\
& - 4m_B\hat{s}(1-\hat{r}_{K^*}-\hat{s})\Big\{(B_4-B_5)B_6^*C_T^* - 2(B_4+B_5)B_6^*C_{TE}^* + \\
& \qquad + 2m_B\hat{m}_\ell(B_3-D_3)B_6^*C_T^*\Big\} + \\
& + 16m_B^2\hat{m}_\ell(1-\hat{r}_{K^*})(1+3\hat{r}_{K^*}-\hat{s})(B_2-D_2)C_T^*T_1^* + \\
& + 8m_B\hat{s}(1+3\hat{r}_{K^*}-\hat{s})\Big\{(B_4-B_5)C_T^*T_1^* - 2(B_4+B_5)C_{TE}^*T_1^*\Big\} - \\
& - 16\hat{m}_\ell(1+3\hat{r}_{K^*}-\hat{s})(B_1-D_1)C_T^*T_1^* + \\
& + 16m_B^2\hat{m}_\ell\hat{s}(1+3\hat{r}_{K^*}-\hat{s})(B_3-D_3)C_T^*T_1^*\Big],
\end{aligned} \tag{3.12}
$$

$$
\begin{aligned}
\mathcal{A}_{FB}^{TT} = \frac{2}{\hat{r}_{K^*}\Delta}m_B^3\sqrt{\lambda}v\,\mathrm{Re}\Big[& m_B^2\hat{m}_\ell\lambda\Big\{4(B_1-D_1)B_7^*C_T^* + (B_2+D_2)(B_4^*+B_5^*)\Big\} - \\
& - 4m_B^4\hat{m}_\ell(1-\hat{r}_{K^*})\lambda(B_2-D_2)B_7^*C_T^* - \\
& - 2m_B^3\hat{s}\lambda\Big\{(B_4-B_5)B_7^*C_T^* - 2(B_4+B_5)B_7^*C_{TE}^* + \\
& \qquad + 2m_B\hat{m}_\ell(B_3-D_3)B_7^*C_T^*\Big\} - \\
& - 2(1-\hat{r}_{K^*}-\hat{s})\Big\{\hat{m}_\ell\Big[B_1^*(B_4+B_5+8B_6C_T) + \\
& \qquad\qquad + D_1^*(B_4+B_5-8B_6C_T)\Big] - \\
& \qquad - 4m_B\hat{s}\Big[(B_4-B_5)B_6^*C_T^* - 2(B_4+B_5)B_6^*C_{TE}^* + \\
& \qquad\qquad + 2m_B\hat{m}_\ell(B_3-D_3)B_6^*C_T^*\Big]\Big\} +
\end{aligned}
$$

$$
\begin{aligned}
&+ 8m_B^2\hat{m}_\ell(1-\hat{r}_{K^*})(1-\hat{r}_{K^*}-\hat{s})(B_2-D_2)B_6^*C_T^* - \\
&- 16m_B^2\hat{m}_\ell(1-\hat{r}_{K^*})(1+3\hat{r}_{K^*}-\hat{s})(B_2-D_2)C_T^*T_1^* - \qquad (3.13)\\
&- 8m_B\hat{s}(1+3\hat{r}_{K^*}-\hat{s})\Big\{(B_4-B_5)C_T^*T_1^* - 2(B_4+B_5)C_{TE}^*T_1^*\Big\} + \\
&+ 16\hat{m}_\ell(1+3\hat{r}_{K^*}-\hat{s})\Big\{(B_1-D_1)C_T^*T_1^* - m_B^2\hat{s}(B_3-D_3)C_T^*T_1^*\Big\}\Big].
\end{aligned}
$$

In these expressions for $\mathcal{A}_{FB}^{ij}$, the first index in the superscript describes the polarization of lepton and the second index describes that of anti-lepton.

It should be noted here that, the double-polarized FB asymmetry for the $B \to K\tau^+\tau^-$ and $b \to s\tau^+\tau^-$ decays are calculated in the supersymmetric model in [23].

4. Numerical analysis

In this section we analyze the effects of the Wilson coefficients on the polarized FB asymmetry. The input parameters we use in our numerical calculations are: $|V_{tb}V_{ts}^*| = 0.0385$, $m_{K^*} = 0.892\,\text{GeV}$, $m_\tau = 1.77\,\text{GeV}$, $m_\mu = 0.106\,\text{GeV}$, $m_b = 4.8\,\text{GeV}$, $m_B = 5.26\,\text{GeV}$ and $\Gamma_B = 4.22\times10^{-13}\,\text{GeV}$. For the values of the Wilson coefficients we use $C_7^{SM} = -0.313$, $C_9^{SM} = 4.344$ and $C_{10}^{SM} = -4.669$. It should be noted that the above-presented value for C_9^{SM} corresponds only to short distance contributions. In addition to the short distance contributions, it receives long distance contributions which result from the conversion of $\bar{c}c$ to the lepton pair. In this work we neglect long distance contributions. The reason for such a choice is dictated by the fact that, in the SM the zero position of $\mathcal{A}_{FB}$ for the $B \to K^*\ell^+\ell^-$ decay is practically independent of the form factors and is determined in terms of short distance Wilson coefficients C_9^{SM} and C_7^{SM} (see [7, 12]) and $s_0 = 3.9\,GeV^2$. For the form factors we have used the light cone QCD sum rules results [24, 25]. As a result of the analysis carried out in this scheme, the q^2 dependence of the form factors can be represented in terms of three parameters as

$$
F(q^2) = \frac{F(0)}{1 - a_F\hat{s} + b_F\hat{s}^2}\,,
$$

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.

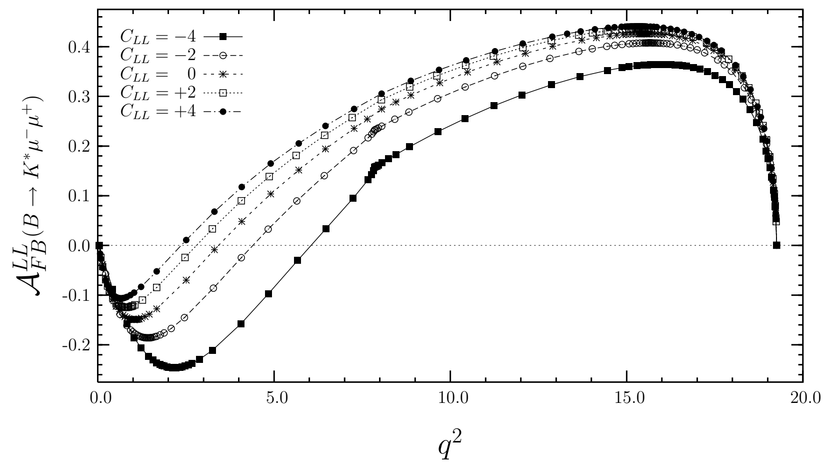

Figure 1: The dependence of the double-lepton polarization asymmetry $\mathcal{A}_{FB}^{LL}$ on q^2 at four fixed values of C_{LL}, for the $B \to K^*\mu^+\mu^-$ decay.

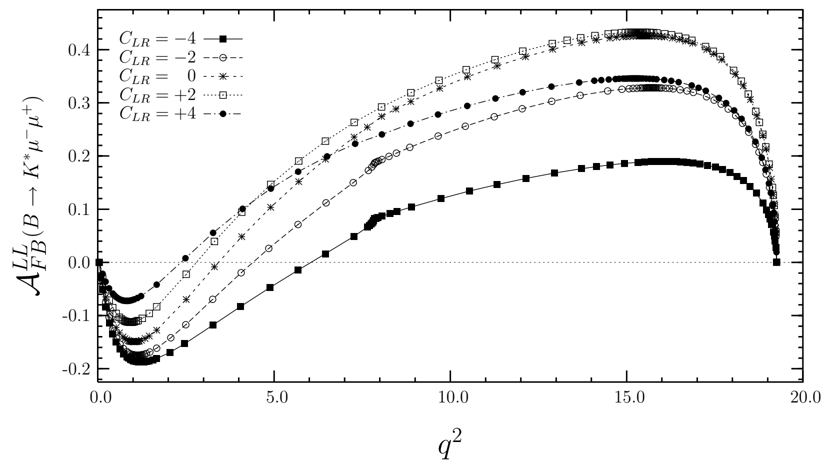

Figure 2: The same as in figure 1, but at four fixed values of C_{LR}.

The new Wilson coefficients vary in the range $-|C_{10}| \leq |C_i| \leq |C_{10}|$. The experimental value of the branching ratio of the $B \to K^*\ell^+\ell^-$ decay [20, 21] and the bound on the branching ratio of the $B \to \mu^+\mu^-$ [26] suggest that this is the right order of magnitude for the vector and scalar interaction coefficients. It should be noted here that the experimental results lead to stronger restrictions on some of the Wilson coefficients, namely $-1.5 \leq C_T \leq 1.5$, $-3.3 \leq C_{TE} \leq 2.6$, $-2 \leq C_{LL}$, $C_{RL} \leq 2.3$, while the remaining coefficients vary in the range $-4 \leq C_X \leq 4$. Since all existing experimental results are yet preliminary, we will vary all new Wilson coefficients in the range $-4 \leq C_X \leq 4$.

In figure 1 (2) we present the dependence of the $\mathcal{A}_{FB}^{LL}$ on q^2 for the $B \to K^*\mu^+\mu^-$ at four fixed values of $C_{LL}(C_{LR})$: $-4, -2, 2, 4$. From these figures we see that nonzero values of the new Wilson coefficients shift the zero position of $\mathcal{A}_{FB}^{LL}$ corresponding to the SM result. When C_{LL} gets negative (positive) values, the zero position of $\mathcal{A}_{FB}^{LL}$ shifts to the left (right) in comparison to that of the zero position in the SM.

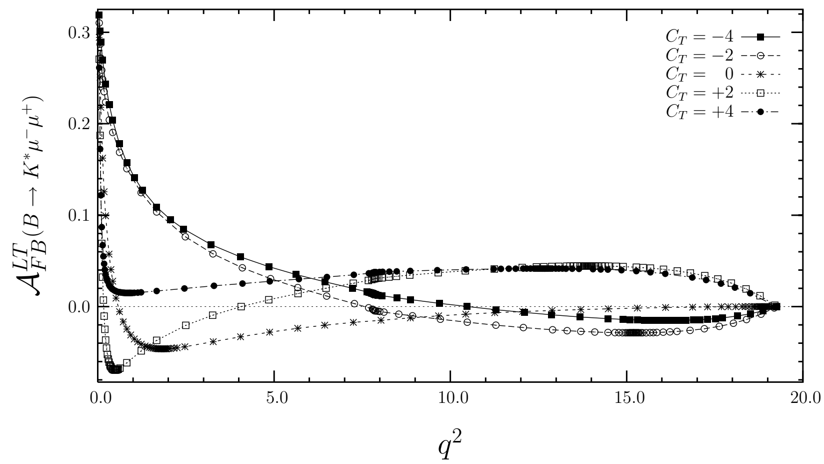

Figure 3: The dependence of the double-lepton polarization asymmetry $\mathcal{A}_{FB}^{LT}$ on q^2 at four fixed values of C_T, for the $B \to K^* \mu^+ \mu^-$ decay.

Our analysis shows that the zero position of $\mathcal{A}_{FB}^{LL}$ for the $B \to K^* \mu^+ \mu^-$ decay is practically independent of the existence of other Wilson coefficients. For this reason we do not present the dependence of $\mathcal{A}_{FB}^{LL}$ on q^2 at fixed values of the remaining Wilson coefficients.

Figures 3 (5) and 4 (6) depict the dependence of $\mathcal{A}_{FB}^{LT}$ and $\mathcal{A}_{FB}^{TL}$ on q^2 at four fixed values of $C_T(C_{TE})$. We observe from these figures that the zero positions of $\mathcal{A}_{FB}^{LT}$ and $\mathcal{A}_{FB}^{TL}$ are very sensitive to the existence of tensor interactions. More essential than is that in the SM case $\mathcal{A}_{FB}^{LT}$ and $\mathcal{A}_{FB}^{TL}$ do not have zero values. Therefore, if zero values for the polarized $\mathcal{A}_{FB}^{LT}$ and $\mathcal{A}_{FB}^{TL}$ asymmetries are measured in the experiments in future, these results are unambiguous indication of the existence of new physics beyond the SM, more specifically, the existence of tensor interactions.

In the case of $B \to K^* \tau^+ \tau^-$ decay, the zero position for the double polarization asymmetries $\mathcal{A}_{FB}^{ij}$ is absent for most of the new Wilson coefficients, and hence, it could be concluded to be insensitive to the new physics beyond the SM, or the value of $\mathcal{A}_{FB}^{ij}$ is quite small, whose measurement in the experiments could practically be impossible. For this reason we do not present the dependencies of $\mathcal{A}_{FB}^{ij}$ on q^2 at fixed values of C_X for the $B \to K^* \tau^+ \tau^-$ decay.

As is obvious from the explicit expressions of the forward-backward asymmetries, they depend both on q^2 and the new Wilson coefficients C_X. As a result of this, it might be difficult to study the dependence of the polarized forward-backward asymmetries $\mathcal{A}_{FB}^{ij}$ on these parameters. However, we can eliminate the dependence of the polarized $\mathcal{A}_{FB}^{ij}$ on q^2 by performing integration over q^2 in the kinematically allowed region, so that the polarized forward-backward asymmetry is said to be averaged. The averaged polarized forward-backward asymmetry is defined as

$$\left\langle \mathcal{A}_{FB}^{ij} \right\rangle = \frac{\int_{4m_\ell^2}^{(m_B - m_{K^*})^2} \mathcal{A}_{FB}^{ij} \frac{d\mathcal{B}}{dq^2} dq^2}{\int_{4m_\ell^2}^{(m_B - m_{K^*})^2} \frac{d\mathcal{B}}{dq^2} dq^2} \,.$$

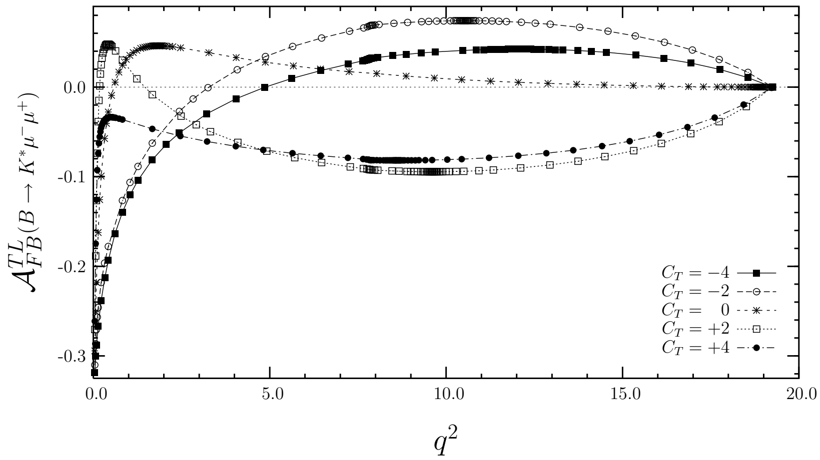

Figure 4: The same as in figure 3, but for $\mathcal{A}_{FB}^{TL}$.

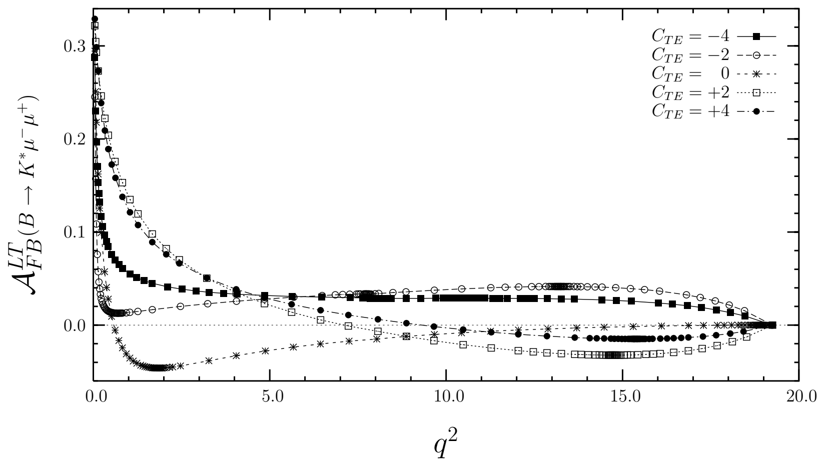

Figure 5: The same as in figure 3, but at four fixed values of C_{TE}.

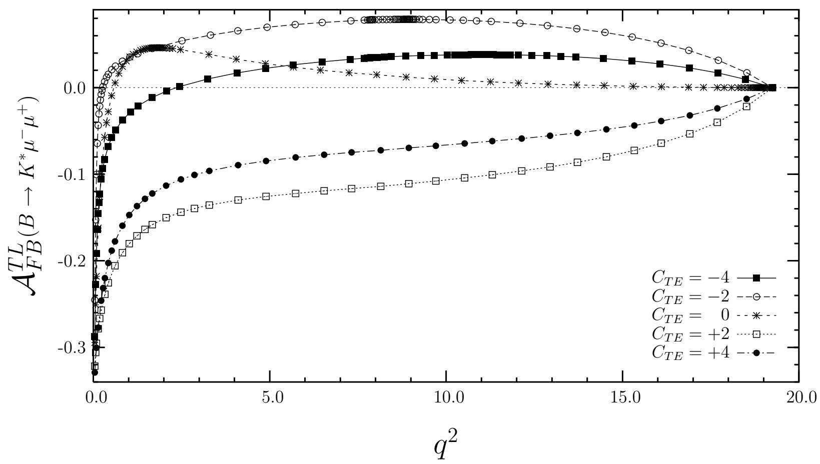

Figure 6: The same as in figure 4, but at four fixed values of C_{TE}.

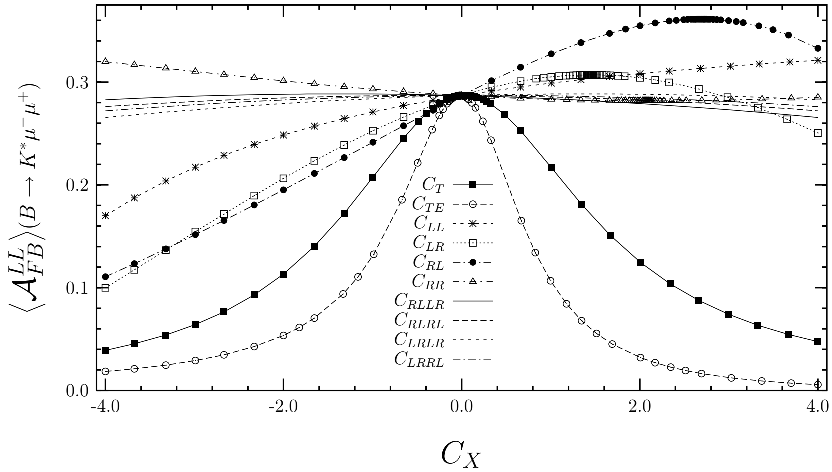

Figure 7: The dependence of the averaged forward-backward double-lepton polarization asymmetry $\langle \mathcal{A}_{FB}^{LL} \rangle$ on the new Wilson coefficients C_X, for the $B \to K^* \mu^+ \mu^-$ decay.

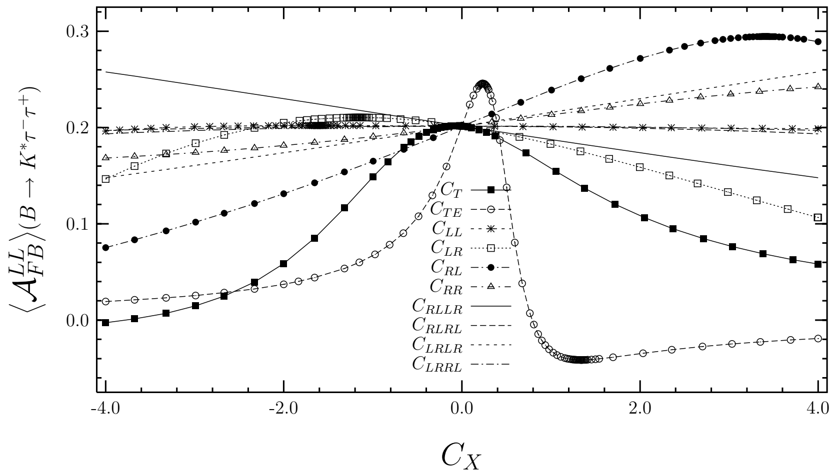

Figure 8: The same as in figure 7, but for the $B \to K^* \tau^+ \tau^-$ decay.

In figure 7, we present the dependence of $\langle \mathcal{A}_{FB}^{LL} \rangle$ on C_X for the $B \to K^* \mu^+ \mu^-$ decay. The common intersection point of all curves corresponds to the SM case. We observe from this figure that, $\langle \mathcal{A}_{FB}^{LL} \rangle$ has symmetric behavior on its dependence on C_T and C_{TE} with respect to zero position, and remains smaller compared to the SM result. The only case for which $\langle \mathcal{A}_{FB}^{LL} \rangle > \langle \mathcal{A}_{FB}^{LL} \rangle_{SM}$ occurs for the positive values of the vector interaction coefficients. Therefore, if we measure in the experiments $\langle \mathcal{A}_{FB}^{LL} \rangle > \langle \mathcal{A}_{FB}^{LL} \rangle_{SM}$, it is a direct indication of new physics beyond the SM, and this departure is to be attributed solely to the existence of vector type interactions.

The situation is even more conformative for the $B \to K^* \tau^+ \tau^-$ case. In figures 9 and 10, we present the dependencies of $\langle \mathcal{A}_{FB}^{LL} \rangle$ and $\langle \mathcal{A}_{FB}^{LT} \rangle$ on the new Wilson coefficients C_X. From figure 9 we observe that, with respect to the zero value of the Wilson coefficients, $\langle \mathcal{A}_{FB}^{LL} \rangle$ increases if C_{RL}, C_{LRLR} and C_{RR} increase, while it decreases when C_{RLLR} increases.

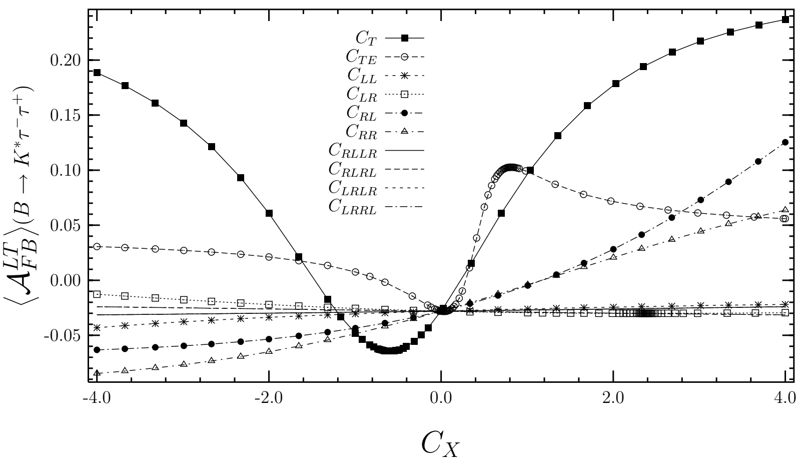

Figure 9: The same as in figure 8, but for the averaged forward-backward double-lepton polarization asymmetry $\langle \mathcal{A}_{FB}^{LT} \rangle$.

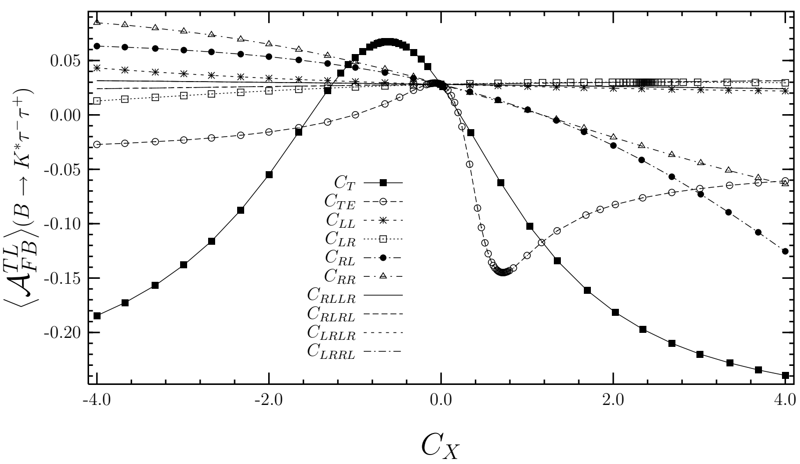

Figure 10: The same as in figure 8, but for the averaged forward-backward double-lepton polarization asymmetry $\langle \mathcal{A}_{FB}^{TL} \rangle$.

From figure 9 we see that the dependence of $\langle \mathcal{A}_{FB}^{LT} \rangle$ on the tensor interaction is stronger. When C_T, C_{TE} and C_{LR} are negative (positive) and vary from -4 to zero (from zero to 4) $\langle \mathcal{A}_{FB}^{LT} \rangle$ decrease (increase). Additionally, we observe that with increasing values of C_{RL} and C_{RR}, $\langle \mathcal{A}_{FB}^{LT} \rangle$ increases. This figure further depicts that $\langle \mathcal{A}_{FB}^{LT} \rangle$, for practical purposes, is not sensitive to the existence of scalar interactions. On the other hand, $\langle \mathcal{A}_{FB}^{NN} \rangle$ and $\langle \mathcal{A}_{FB}^{TT} \rangle$ are very sensitive to the presence of tensor and scalar interactions (see figures 10 and 11).

It is clear from these results that several of the polarized forward-backward asymmetries show sizable departure from the SM results and they are sensitive to the existence of different type of interactions. therefore, study of these observables can be very useful in looking for new physics beyond the SM.

Obviously, if new physics beyond the SM exists, there hoped to be effects on the branching ratio besides the polarized $\mathcal{A}_{FB}$. Keeping in mind that the measurement of the branching ratio is easier, one could find it more convenient to study it for establishing new physics. But the intriguing question is, whether there could appear situations in which

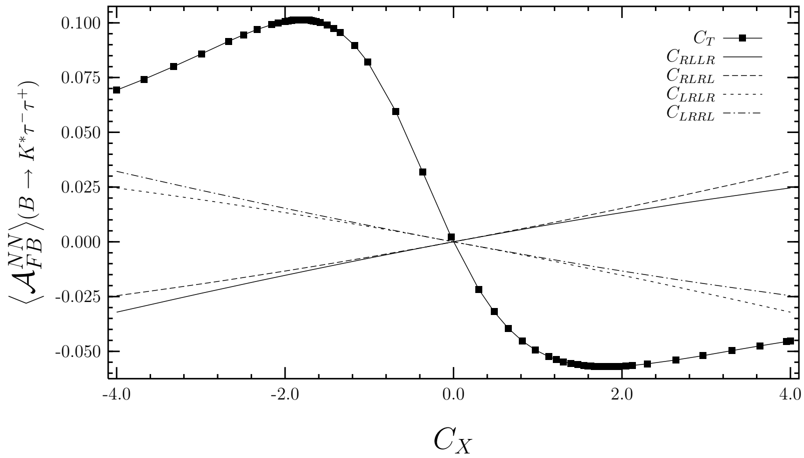

Figure 11: The same as in figure 8, but for the averaged forward-backward double-lepton polarization asymmetry $\langle \mathcal{A}_{FB}^{NN} \rangle$.

the value of the branching ratio coincides with that of the SM result, while polarized $\mathcal{A}_{FB}$ does not. In order to answer this question we study the correlation between the averaged, polarized $\langle \mathcal{A}_{FB} \rangle$ and branching ratio. In further analysis we vary the branching ratio of $B \to K^* \mu^+ \mu^-$ ($K^* \tau^+ \tau^-$) between the values $(1-3) \times 10^{-6}$ $[(1-3) \times 10^{-7}]$, which is very close to the SM calculations. Note that, we do not take into account the experimental results on branching ratio since they contain large errors, and it would be better to wait for more improved experimental results.

Our conclusion for the $B \to K^* \mu^+ \mu^-$ decay, in regard to the above-mentioned correlated relation, is as follows (remember that, the intersection of all curves corresponds to the SM value):

- for $\langle \mathcal{A}_{FB}^{LL} \rangle$, such a region is absent for all C_X,
- for $\langle \mathcal{A}_{FB}^{TL} \rangle$, such a region does exist for C_T and C_{TE} (see figure 13).

The situation is much more attractive for the $B \to K^* \tau^+ \tau^-$ decay. In figures 14–18, we depict the dependence of the averaged, forward-backward polarized asymmetries $\langle \mathcal{A}_{FB}^{LL} \rangle$; $\langle \mathcal{A}_{FB}^{LT} \rangle \approx -\langle \mathcal{A}_{FB}^{TL} \rangle$; $\langle \mathcal{A}_{FB}^{NT} \rangle \approx \langle \mathcal{A}_{FB}^{TN} \rangle$; $\langle \mathcal{A}_{FB}^{NN} \rangle$ and $\langle \mathcal{A}_{FB}^{TT} \rangle$, on branching ratio. It follows from these figures that, indeed, there exist certain regions of the new Wilson coefficients for which, mere study of the polarized $\mathcal{A}_{FB}$ can give promising information about new physics beyond the SM.

In summary, in this work we present the analysis for the forward-backward asymmetries when both leptons are polarized, using a general, model independent form of the effective hamiltonian. Our work verifies that the study of the zero position of $\langle \mathcal{A}_{FB}^{LL} \rangle$ can give unambiguous conformation of the new physics beyond the SM, since when new physics effects are taken into account, the results are shifted with respect to their zero positions in the SM. Moreover, we find that the polarized $\mathcal{A}_{FB}$ is quite sensitive to the existence of the tensor and vector interactions. Finally we obtain that there exist certain regions of the new Wilson coefficients for which, only study of the polarized forward-backward asymmetry gives invaluable information in establishing new physics beyond the SM.

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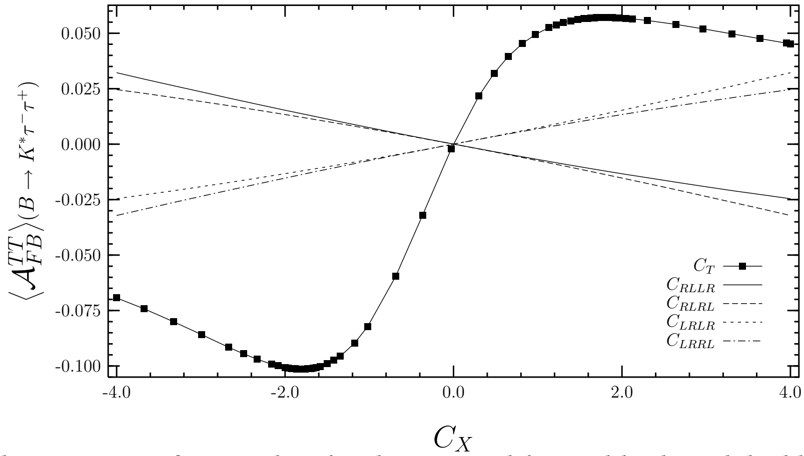

Figure 12: The same as in figure 8, but for the averaged forward-backward double-lepton polarization asymmetry $\langle \mathcal{A}_{FB}^{TT} \rangle$.

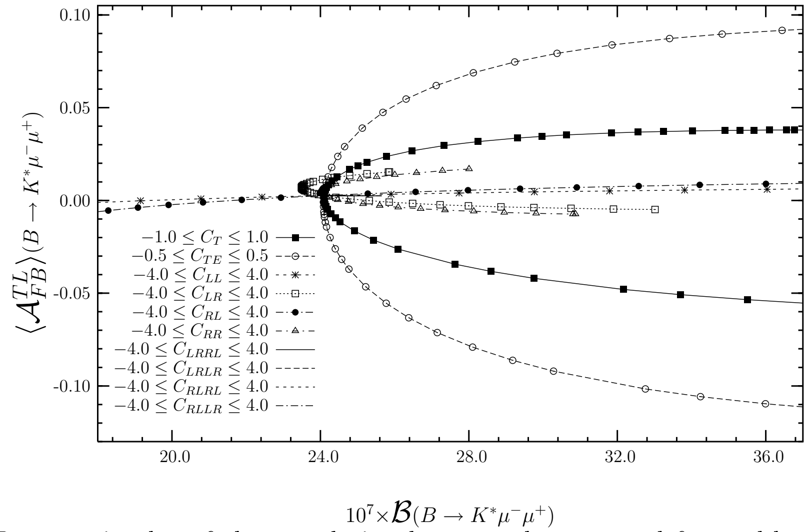

Figure 13: Parametric plot of the correlation between the averaged forward-backward double-lepton polarization asymmetry $\langle \mathcal{A}_{FB}^{LT} \rangle$ and the branching ratio for the $B \to K^* \mu^+ \mu^-$ decay.

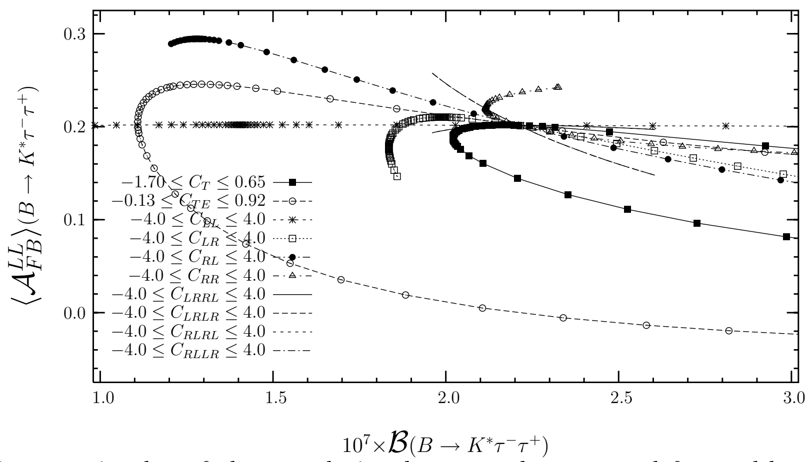

Figure 14: Parametric plot of the correlation between the averaged forward-backward double-lepton polarization asymmetry $\langle \mathcal{A}_{FB}^{LL} \rangle$ and the branching ratio for the $B \to K^* \tau^+ \tau^-$ decay.

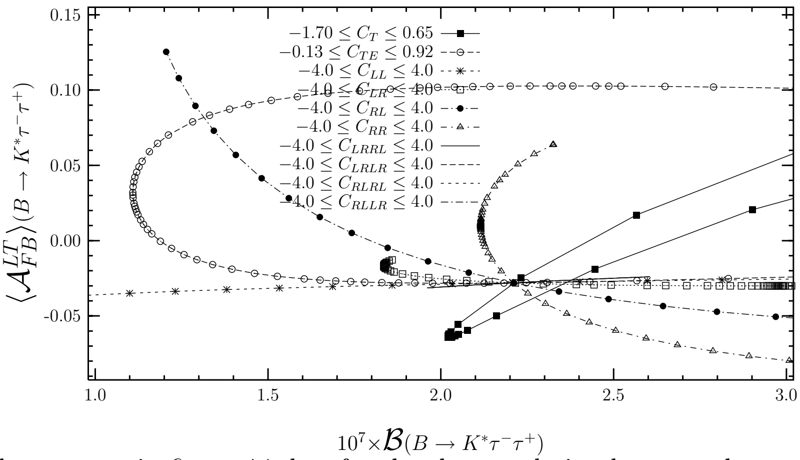

Figure 15: The same as in figure 14, but for the the correlation between the averaged forward-backward double-lepton polarization asymmetry $\langle \mathcal{A}_{FB}^{LT} \rangle$ and the branching ratio.

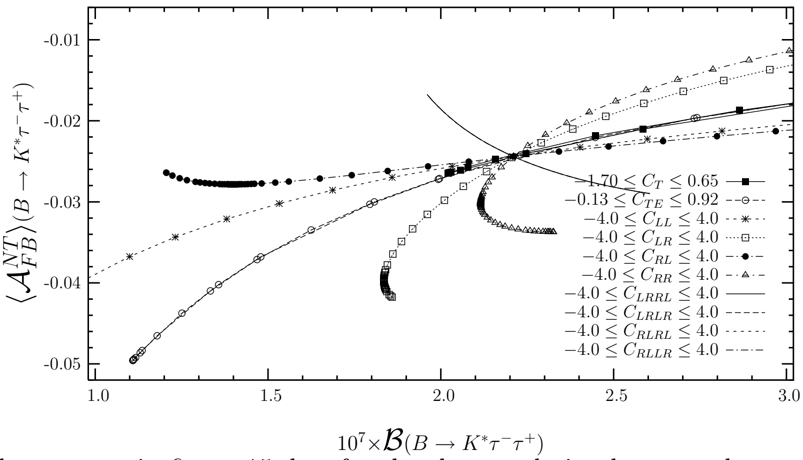

Figure 16: The same as in figure 15, but for the the correlation between the averaged forward-backward double-lepton polarization asymmetry $\langle \mathcal{A}_{FB}^{NT} \rangle$ and the branching ratio.

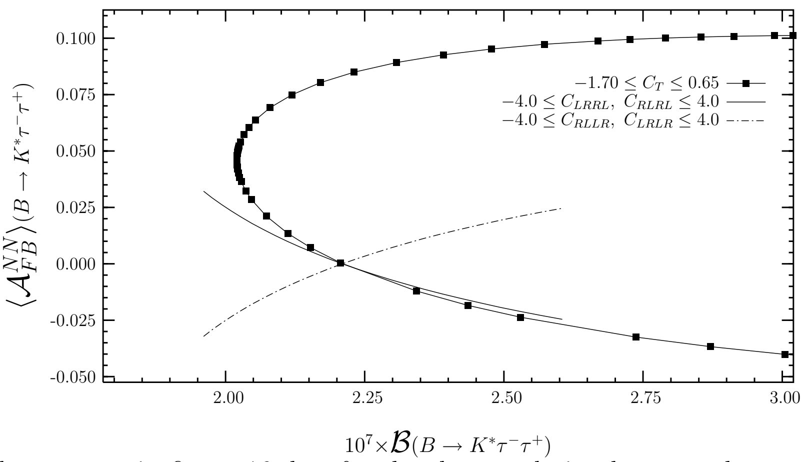

Figure 17: The same as in figure 16, but for the the correlation between the averaged forward-backward double-lepton polarization asymmetry $\langle \mathcal{A}_{FB}^{NN} \rangle$ and the branching ratio.

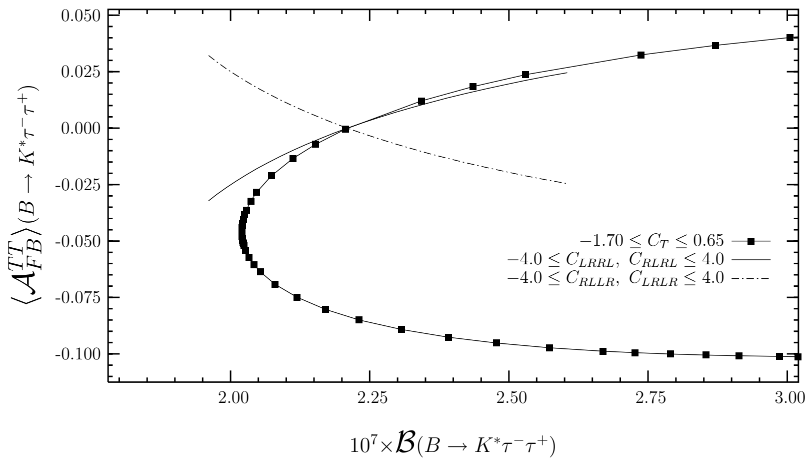

Figure 18: The same as in figure 17, but for the the correlation between the averaged forward-backward double-lepton polarization asymmetry $\langle \mathcal{A}_{FB}^{TT} \rangle$ and the branching ratio.

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