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Polarization effects in exclusive semileptonic $\Lambda_b \to \Lambda \ell^+ \ell^-$ decay

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Abstract: The independent helicity amplitudes in the $\Lambda_b \to \Lambda \ell^+ \ell^-$ decay in the standard model and its minimal extension, i.e., with the new vector type interactions, are calculated. We calculate various asymmetry parameters characterizing the angular dependence of the differential decay width for the cascade decay $\Lambda_b \to \Lambda(\to a+b)\, V^*(\to \ell^+ \ell^-)$ with polarized and unpolarized heavy baryons. The sensitivity of the asymmetry parameters to the new Wilson coefficients are analyzed.

Keywords: Beyond Standard Model, QCD.

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Contents

1. Introduction

Rare B-decays induced by the flavor-changing neutral current (FCNC) $b \to s$ or $b \to d$ transitions occur at loop level in the standard model (SM), since FCNC transitions that are forbidden in the SM at tree level provide consistency check of the SM at quantum level. These decays induced by the FCNC are also very promising tools for establishing new physics beyond the SM. New physics appear in rare decays through the Wilson coefficients which can take values different from their SM counterpart or through the new operator structures in an effective Hamiltonian (see for example [1]–[13]).

Among the hadronic, leptonic and semileptonic decays, the last decay channels are very significant, since they are theoretically, more or less, clean, and they have relatively larger branching ratio. The semileptonic decay channels is described by the $b \to s(d)\ell^+\ell^-$ transition and they contain many observables like forward-backward asymmetry $\mathcal{A}_{FB}$, lepton polarization asymmetries, etc. Existence of these observables is very useful and serve as a testing ground for the SM and in looking for new physics beyond th SM. For this reason, many processes, like $B \to \pi(\rho)\ell^+\ell^-$ [14], $B \to \ell^+\ell^-\gamma$ [15], $B \to K\ell^+\ell^-$ [16] and $B \to K^*\ell^+\ell^-$ [17–24] have been studied comprehensively.

Recently, BELLE and BaBar Collaborations announced the following results for the branching ratios of the $B \to K^*\ell^+\ell^-$ and $B \to K\ell^+\ell^-$ decays:

$$\mathcal{B}(B \to K^*\ell^+\ell^-) = \begin{cases} \left(11.5^{+2.6}_{-2.4} \pm 0.8 \pm 0.2\right) \times 10^{-7} & [25]\,, \\ \left(0.78^{+0.19}_{-0.17} \pm 0.12\right) \times 10^{-6} & [26]\,, \end{cases}$$

$$\mathcal{B}(B \to K\ell^+\ell^-) = \begin{cases} \left(4.8^{+1.0}_{-0.9} \pm 0.3 \pm 0.1\right) \times 10^{-7} & [25]\,, \\ (0.34 \pm 0.07 \pm 0.12) \times 10^{-6} & [26]\,. \end{cases}$$

Another exclusive decay which is described at inclusive level by the $b \to s\ell^+\ell^-$ transition is the baryonic $\Lambda_b \to \Lambda\ell^+\ell^-$ decay. Unlike mesonic decays, the baryonic decays could maintain the helicity structure of the effective Hamiltonian for the $b \to s$ transition [27]. Radiative and semileptonic decays of Λ_b such as $\Lambda_b \to \Lambda\gamma$, $\Lambda_b \to \Lambda_c\ell\bar{\nu}_\ell$, $\Lambda_b \to \Lambda\ell^+\ell^-$ ($\ell = e, \mu, \tau$) and $\Lambda_b \to \Lambda\nu\bar{\nu}$ have been extensively studied in the literature [28–33]. More details about heavy baryons, including the experimental prospects, can be found in [34, 35].

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Many experimentally measurable quantities such as branching ratio [36], Λ polarization and single- and double-lepton polarizations have already been studied in [37, 38] and [39], respectively. Analysis of such quantities can be useful for more precise determination of the SM parameters and in looking for new physics beyond the SM.

In the present work we analyze the possibility of searching for new physics in the baryonic $\Lambda_b \to \Lambda\ell^+\ell^-$ decay by studying different asymmetry parameters that characterize the angular dependence of the angular decay distributions, with the inclusion of non-standard vector type of interactions. In our analysis we use the helicity amplitude formalism and polarization density matrix method (see the first and third references in [28]) to analyze the joint decay distributions in this decay.

The paper is organized as follows. In section 2, using the Hamiltonian that includes non-standard vector interactions, the matrix element for the $\Lambda_b \to \Lambda\ell^+\ell^-$ is obtained. In section 3 we calculate the different polarization asymmetries. In the final section we study the sensitivity of various asymmetries to the non-standard interactions.

2. Matrix element for the $\Lambda_b \to \Lambda\ell^+\ell^-$ decay

In this section we derive the matrix element for the $\Lambda_b \to \Lambda\ell^+\ell^-$ decay which is described by the $b \to s\ell^+\ell^-$ transition at quark level. Neglecting the terms proportional to $V_{ub}V_{us}^*/V_{tb}V_{ts}^* \sim \mathcal{O}(10^{-2})$, the matrix element for the $b \to s\ell^+\ell^-$ decay can be written in terms of the twelve model independent four-Fermi interactions as [18]

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

where $q = p_{\Lambda_b} - p_\Lambda = p_1 + p_2$ is the momentum transfer and C_X are the coefficients of the four-Fermi interactions, $L = (1-\gamma_5)/2$ and $R = (1+\gamma_5)/2$. The terms with coefficients C_{SL} and C_{BR} describe the penguin contributions, which correspond to $-2m_sC_7^{\rm eff}$ and $-2m_bC_7^{\rm eff}$ in the SM, respectively. The next four terms in eq. (2.1) with coefficients $C_{LL}^{\rm tot}$, $C_{LR}^{\rm tot}$, C_{RL} and C_{RR} describe vector type interactions, two ($C_{LL}^{\rm tot}$ and $C_{LR}^{\rm tot}$) of which contain SM contributions in the form $C_9^{\rm eff} - C_{10}$ and $C_9^{\rm eff} + C_{10}$, respectively. Thus, $C_{LL}^{\rm tot}$ and $C_{LR}^{\rm tot}$ can be written as

$$\begin{aligned} C_{LL}^{\rm tot} &= C_9^{\rm eff} - C_{10} + C_{LL}\,, \\ C_{LR}^{\rm tot} &= C_9^{\rm eff} + C_{10} + C_{LR}\,, \end{aligned} \tag{2.2}$$

where C_{LL} and C_{LR} describe the contributions of new physics. Additionally, eq. (2.1) contains four scalar type interactions (C_{LRLR}, C_{RLLR}, C_{LRRL} and C_{RLRL}), and two tensor type interactions (C_T and C_{TE}). In the present work we will consider the minimal extension of the SM and therefore we neglect the scalar and tensor type interactions throughout in this work.

The amplitude of the exclusive $\Lambda_b \to \Lambda \ell^+ \ell^-$ decay is obtained by calculating the matrix element of $\mathcal{M}$ for the $b \to s \ell^+ \ell^-$ transition between initial and final baryon states $\langle \Lambda | \mathcal{M} | \Lambda_b \rangle$. It follows from eq. (2.1) that the matrix elements

$$\begin{aligned} &\langle \Lambda | \bar{s} \gamma_\mu (1 \mp \gamma_5) b | \Lambda_b \rangle \ , \\ &\langle \Lambda | \bar{s} \sigma_{\mu\nu} (1 \mp \gamma_5) b | \Lambda_b \rangle \ , \end{aligned}$$

are needed in order to calculate the $\Lambda_b \to \Lambda \ell^+ \ell^-$ decay amplitude.

These matrix elements parametrized in terms of the form factors are as follows (see [37, 40])

$$\langle \Lambda | \bar{s} \gamma_\mu b | \Lambda_b \rangle = \bar{u}_\Lambda \Big[f_1 \gamma_\mu + i f_2 \sigma_{\mu\nu} q^\nu + f_3 q_\mu \Big] u_{\Lambda_b} \ , \tag{2.3}$$

$$\langle \Lambda | \bar{s} \gamma_\mu \gamma_5 b | \Lambda_b \rangle = \bar{u}_\Lambda \Big[g_1 \gamma_\mu \gamma_5 + i g_2 \sigma_{\mu\nu} \gamma_5 q^\nu + g_3 q_\mu \gamma_5 \Big] u_{\Lambda_b} \ , \tag{2.4}$$

$$\langle \Lambda | \bar{s} \sigma_{\mu\nu} b | \Lambda_b \rangle = \bar{u}_\Lambda \Big[f_T \sigma_{\mu\nu} - i f_T^V (\gamma_\mu q^\nu - \gamma_\nu q^\mu) - i f_T^S (P_\mu q^\nu - P_\nu q^\mu) \Big] u_{\Lambda_b} \ , \tag{2.5}$$

$$\langle \Lambda | \bar{s} \sigma_{\mu\nu} \gamma_5 b | \Lambda_b \rangle = \bar{u}_\Lambda \Big[g_T \sigma_{\mu\nu} - i g_T^V (\gamma_\mu q^\nu - \gamma_\nu q^\mu) - i g_T^S (P_\mu q^\nu - P_\nu q^\mu) \Big] \gamma_5 u_{\Lambda_b} \ , \tag{2.6}$$

where $P = p_{\Lambda_b} + p_\Lambda$ and $q = p_{\Lambda_b} - p_\Lambda$.

The form factors of the magnetic dipole operators are defined as

$$\begin{aligned} \langle \Lambda | \bar{s} i \sigma_{\mu\nu} q^\nu b | \Lambda_b \rangle &= \bar{u}_\Lambda \Big[f_1^T \gamma_\mu + i f_2^T \sigma_{\mu\nu} q^\nu + f_3^T q_\mu \Big] u_{\Lambda_b} \ , \\ \langle \Lambda | \bar{s} i \sigma_{\mu\nu} \gamma_5 q^\nu b | \Lambda_b \rangle &= \bar{u}_\Lambda \Big[g_1^T \gamma_\mu \gamma_5 + i g_2^T \sigma_{\mu\nu} \gamma_5 q^\nu + g_3^T q_\mu \gamma_5 \Big] u_{\Lambda_b} \ . \end{aligned} \tag{2.7}$$

Using the identity

$$\sigma_{\mu\nu} \gamma_5 = -\frac{i}{2} \epsilon_{\mu\nu\alpha\beta} \sigma^{\alpha\beta} \ ,$$

and eq. (2.5), the last expression in eq. (2.7) can be written as

$$\langle \Lambda | \bar{s} i \sigma_{\mu\nu} \gamma_5 q^\nu b | \Lambda_b \rangle = \bar{u}_\Lambda \Big[f_T i \sigma_{\mu\nu} \gamma_5 q^\nu \Big] u_{\Lambda_b} \ .$$

Multiplying (2.5) and (2.6) by $i q^\nu$ and comparing with (2.7), one can easily obtain the following relations between the form factors

$$\begin{aligned} f_2^T &= f_T + f_T^S q^2 \ , \\ f_1^T &= \Big[f_T^V + f_T^S (m_{\Lambda_b} + m_\Lambda) \Big] q^2 \ = -\frac{q^2}{m_{\Lambda_b} - m_\Lambda} f_3^T \ , \\ g_2^T &= g_T + g_T^S q^2 \ , \\ g_1^T &= \Big[g_T^V - g_T^S (m_{\Lambda_b} - m_\Lambda) \Big] q^2 = \frac{q^2}{m_{\Lambda_b} + m_\Lambda} g_3^T \ . \end{aligned} \tag{2.8}$$

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Using these definitions of the form factors, for the matrix element of the $\Lambda_b \to \Lambda \ell^+ \ell^-$ we get [37, 38]

$$
\begin{aligned}
\mathcal{M} = \frac{G\alpha}{4\sqrt{2}\pi} V_{tb} V_{ts}^* \frac{1}{2} \Bigg\{ & \bar{\ell}\gamma_\mu (1-\gamma_5)\ell \, \bar{u}_\Lambda \Big[(A_1 - D_1)\gamma_\mu (1+\gamma_5) + (B_1 - E_1)\gamma_\mu (1-\gamma_5) \\
& + i\sigma_{\mu\nu} q^\nu \Big((A_2 - D_2)(1+\gamma_5) + (B_2 - E_2)(1-\gamma_5) \Big) \\
& + q_\mu \Big((A_3 - D_3)(1+\gamma_5) + (B_3 - E_3)(1-\gamma_5) \Big) \Big] u_{\Lambda_b} \\
& + \bar{\ell}\gamma_\mu (1+\gamma_5)\ell \, \bar{u}_\Lambda \Big[(A_1 + D_1)\gamma_\mu (1+\gamma_5) + (B_1 + E_1)\gamma_\mu (1-\gamma_5) \\
& + i\sigma_{\mu\nu} q^\nu \Big((A_2 + D_2)(1+\gamma_5) + (B_2 + E_2)(1-\gamma_5) \Big) \\
& + q_\mu \Big((A_3 + D_3)(1+\gamma_5) + (B_3 + E_3)(1-\gamma_5) \Big) \Big] u_{\Lambda_b} \Bigg\} ,
\end{aligned} \tag{2.9}
$$

where

$$
\begin{aligned}
A_1 &= \frac{1}{q^2} \left(f_1^T - g_1^T \right) C_{SL} + \frac{1}{q^2} \left(f_1^T + g_1^T \right) C_{BR} + \frac{1}{2} \left(f_1 - g_1 \right) \left(C_{LL}^{\rm tot} + C_{LR}^{\rm tot} \right) \\
&\quad + \frac{1}{2} \left(f_1 + g_1 \right) \left(C_{RL} + C_{RR} \right) , \\
A_2 &= A_1 \left(1 \to 2 \right) , \\
A_3 &= A_1 \left(1 \to 3 \right) , \\
B_1 &= A_1 \left(g_1 \to -g_1 ;\; g_1^T \to -g_1^T \right) , \\
B_2 &= B_1 \left(1 \to 2 \right) , \\
B_3 &= B_1 \left(1 \to 3 \right) , \\
D_1 &= \frac{1}{2} \left(C_{RR} - C_{RL} \right) \left(f_1 + g_1 \right) + \frac{1}{2} \left(C_{LR}^{\rm tot} - C_{LL}^{\rm tot} \right) \left(f_1 - g_1 \right) , \\
D_2 &= D_1 \left(1 \to 2 \right) , \\
D_3 &= D_1 \left(1 \to 3 \right) , \\
E_1 &= D_1 \left(g_1 \to -g_1 \right) , \\
E_2 &= E_1 \left(1 \to 2 \right) , \\
E_3 &= E_1 \left(1 \to 3 \right) .
\end{aligned} \tag{2.10}
$$

From these expressions it follows that $\Lambda_b \to \Lambda \ell^+ \ell^-$ decay is described in terms of many form factors. It is shown in [41] that Heavy Quark Effective Theory reduces the number of independent form factors to two (F_1 and F_2) irrelevant of the Dirac structure of the corresponding operators, i.e.,

$$
\left\langle \Lambda(p_\Lambda) \left| \bar{s} \Gamma b \right| \Lambda(p_{\Lambda_b}) \right\rangle = \bar{u}_\Lambda \Big[F_1(q^2) + \not{v} F_2(q^2) \Big] \Gamma u_{\Lambda_b} , \tag{2.11}
$$

where Γ is an arbitrary Dirac structure and $v^\mu = p^\mu_{\Lambda_b}/m_{\Lambda_b}$ is the four-velocity of Λ_b. Comparing the general form of the form factors given in eqs. (2.4)–(2.8) with (2.11), one

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can easily obtain the following relations among them (see also [37, 38, 40])

$$
\begin{aligned}
g_1 &= f_1 = f_2^T = g_2^T = F_1 + \sqrt{\hat{r}_\Lambda} F_2 \,, \\
g_2 &= f_2 = g_3 = f_3 = g_T^V = f_T^V = \frac{F_2}{m_{\Lambda_b}} \,, \\
g_T^S &= f_T^S = 0 \,, \\
g_1^T &= f_1^T = \frac{F_2}{m_{\Lambda_b}} q^2 \,, \\
g_3^T &= \frac{F_2}{m_{\Lambda_b}} \left(m_{\Lambda_b} + m_\Lambda \right) \,, \\
f_3^T &= -\frac{F_2}{m_{\Lambda_b}} \left(m_{\Lambda_b} - m_\Lambda \right) \,,
\end{aligned}
\tag{2.12}
$$

where $\hat{r}_\Lambda = m_\Lambda^2 / m_{\Lambda_b}^2$.

In order to obtain the helicity amplitudes for the $\Lambda_b \to \Lambda \ell^+ \ell^-$ decay, it is convenient to regard this decay as a quasi two-body decay $\Lambda_b \to \Lambda V^*$ followed by the leptonic decay $V^* \to \ell^+ \ell^-$, where V^* is the off-shell γ or Z bosons. The matrix element of $\Lambda_b \to \Lambda \ell^+ \ell^-$ decay can be written in the following form:

$$
\mathcal{M}_{\lambda_i}^{\lambda_\ell \bar{\lambda}_\ell} = \sum_{\lambda_{V^*}} \eta_{\lambda_{V^*}} L_{\lambda_{V^*}}^{\lambda_\ell \bar{\lambda}_\ell} H_{\lambda_{V^*}}^{\lambda_i} \,,
$$

where

$$
L_{\lambda_{V^*}}^{\lambda_\ell \bar{\lambda}_\ell} = \varepsilon_{V^*}^{\mu} \left\langle \ell^-(p_\ell, \lambda_\ell) \, \ell^+(p_\ell, \bar{\lambda}_\ell) \left| J_\mu^\ell \right| 0 \right\rangle \,, \tag{2.13}
$$

$$
H_{\lambda_{V^*}}^{\lambda_i} = \left(\varepsilon_{V^*}^{\mu} \right)^* \left\langle \Lambda(p_\Lambda, \lambda_\Lambda) \left| J_\mu^i \right| \Lambda_b(p_{\Lambda_b}) \right\rangle \,, \tag{2.14}
$$

where $\varepsilon_{V^*}^{\mu}$ is the polarization vector of the virtual intermediate vector boson. The metric tensor can be expressed in terms of the polarization vector of the virtual vector particle $\varepsilon_V = \varepsilon(\lambda_V)$ as

$$
-g^{\mu\nu} = \sum_{\lambda_{V^*}} \eta_{\lambda_{V^*}} \varepsilon_{\lambda_{V^*}}^{\mu} \varepsilon_{\lambda_{V^*}}^{*\nu} \,,
$$

where the summation is over the helicity of the virtual vector particle V, $\Lambda_V = \pm 1, 0, t$ with the metric $\eta_\pm = \eta_0 = -\eta_t = 1$, where $\lambda_V = t$ is the scalar (zero) helicity component of the virtual V particle (for more details see [42, 43] and first and third references in [28]). The upper indices in eqs. (2.13) and (2.14) correspond to the helicities of the leptons and the lower ones correspond to the helicity of the Λ baryon. Moreover, J_μ^ℓ and J_μ^i in eqs. (2.13) and (2.14) are the leptonic and hadronic currents, respectively.

In the calculations of the leptonic and baryonic amplitudes we will use two different frames. The leptonic amplitude $L_{\lambda_{V^*}}^{\lambda_\ell \bar{\lambda}_\ell}$ is calculated in the rest frame of the virtual vector boson wit the z-axis chosen along the Λ direction and the x-z plane chosen as the virtual V decay plane. The hadronic amplitude is calculated in the rest frame of Λ_b baryon.

Using eqs. (2.9)–(2.14), after lengthy calculations, we get for the helicity amplitudes:

$$
\begin{aligned}
\mathcal{M}^{++}_{+1/2} &= 2m_\ell \sin\theta\Big(H^{(1)}_{+1/2,+1} + H^{(2)}_{+1/2,+1}\Big) + 2m_\ell\cos\theta\Big(H^{(1)}_{+1/2,0} + H^{(2)}_{+1/2,0}\Big) \\
&\quad + 2m_\ell\Big(H^{(1)}_{+1/2,t} - H^{(2)}_{+1/2,t}\Big)\,, \\
\mathcal{M}^{+-}_{+1/2} &= -\sqrt{q^2}(1-\cos\theta)\Big[(1-v)H^{(1)}_{+1/2,+1} + (1+v)H^{(2)}_{+1/2,+1}\Big] \\
&\quad -\sqrt{q^2}\sin\theta\Big[(1-v)H^{(1)}_{+1/2,0} + (1+v)H^{(2)}_{+1/2,0}\Big]\,, \\
\mathcal{M}^{-+}_{+1/2} &= \sqrt{q^2}(1+\cos\theta)\Big[(1+v)H^{(1)}_{+1/2,+1} + (1-v)H^{(2)}_{+1/2,+1}\Big] \\
&\quad -\sqrt{q^2}\sin\theta\Big[(1+v)H^{(1)}_{+1/2,0} + (1-v)H^{(2)}_{+1/2,0}\Big]\,, \\
\mathcal{M}^{--}_{+1/2} &= -2m_\ell \sin\theta\Big(H^{(1)}_{+1/2,+1} + H^{(2)}_{+1/2,+1}\Big) - 2m_\ell\cos\theta\Big(H^{(1)}_{+1/2,0} + H^{(2)}_{+1/2,0}\Big) \\
&\quad + 2m_\ell\Big(H^{(1)}_{+1/2,t} - H^{(2)}_{+1/2,t}\Big)\,, \\
\mathcal{M}^{++}_{-1/2} &= -2m_\ell \sin\theta\Big(H^{(1)}_{-1/2,-1} + H^{(2)}_{-1/2,-1}\Big) + 2m_\ell\cos\theta\Big(H^{(1)}_{-1/2,0} + H^{(2)}_{-1/2,0}\Big) \\
&\quad + 2m_\ell\Big(H^{(1)}_{-1/2,t} - H^{(2)}_{-1/2,t}\Big)\,, \\
\mathcal{M}^{+-}_{-1/2} &= -\sqrt{q^2}(1+\cos\theta)\Big[(1-v)H^{(1)}_{-1/2,-1} + (1+v)H^{(2)}_{-1/2,-1}\Big] \\
&\quad -\sqrt{q^2}\sin\theta\Big[(1-v)H^{(1)}_{-1/2,0} + (1+v)H^{(2)}_{-1/2,0}\Big]\,, \\
\mathcal{M}^{-+}_{-1/2} &= \sqrt{q^2}(1-\cos\theta)\Big[(1+v)H^{(1)}_{-1/2,-1} + (1-v)H^{(2)}_{-1/2,-1}\Big] \\
&\quad -\sqrt{q^2}\sin\theta\Big[(1+v)H^{(1)}_{-1/2,0} + (1-v)H^{(2)}_{-1/2,0}\Big]\,, \\
\mathcal{M}^{--}_{-1/2} &= 2m_\ell \sin\theta\Big(H^{(1)}_{-1/2,-1} + H^{(2)}_{-1/2,-1}\Big) - 2m_\ell\cos\theta\Big(H^{(1)}_{-1/2,0} + H^{(2)}_{-1/2,0}\Big) \\
&\quad + 2m_\ell\Big(H^{(1)}_{-1/2,t} - H^{(2)}_{-1/2,t}\Big)\,, \qquad (2.15)
\end{aligned}
$$

where

$$
\begin{aligned}
H^{(1)}_{\pm1/2,\pm1} &= H^{(1)\,V}_{1/2,1} \pm H^{(1)\,A}_{1/2,1}\,, \\
H^{(2)}_{\pm1/2,\pm1} &= H^{(2)\,V}_{1/2,1} \pm H^{(2)\,A}_{1/2,1}\,, \\
H^{(1,2)}_{\pm1/2,0} &= H^{(1,2)\,V}_{1/2,0} \pm H^{(1,2)\,A}_{1/2,1}\,, \\
H^{(1,2)}_{\pm1/2,t} &= H^{(1,2)\,V}_{1/2,t} \pm H^{(1,2)\,A}_{1/2,t}\,, \qquad (2.16)
\end{aligned}
$$

where θ is the angle of the positron in the rest frame of the intermediate boson with respect to its helicity axes. Explicit expressions of the helicity amplitudes $H^{V,A}_{\lambda,\lambda_W}$ are

$$
\begin{aligned}
H^{(1)\,V}_{1/2,1} &= -\sqrt{Q_-}\Big[F_1^V - (m_{\Lambda_b} + m_\Lambda)F_2^V\Big]\,, \\
H^{(1)\,A}_{1/2,1} &= -\sqrt{Q_+}\Big[F_1^A + (m_{\Lambda_b} - m_\Lambda)F_2^A\Big]\,, \\
H^{(2)\,V}_{1/2,1} &= H^{(1)\,V}_{1/2,1}(F_1^V \to F_3^V,\ F_2^V \to F_4^V)\,, \\
H^{(2)\,A}_{1/2,1} &= H^{(1)\,A}_{1/2,1}(F_1^A \to F_3^A,\ F_2^A \to F_4^A)\,,
\end{aligned}
$$

$$
\begin{aligned}
H^{(1)\,V}_{1/2,0} &= -\frac{1}{\sqrt{q^2}}\Big\{\sqrt{Q_-}\Big[(m_{\Lambda_b}+m_\Lambda)F_1^V - q^2F_2^V\Big]\Big\}\,,\\
H^{(1)\,A}_{1/2,0} &= -\frac{1}{\sqrt{q^2}}\Big\{\sqrt{Q_+}\Big[(m_{\Lambda_b}-m_\Lambda)F_1^A + q^2F_2^A\Big]\Big\}\,,\\
H^{(2)\,V}_{1/2,0} &= H^{(1)\,V}_{1/2,0}(F_1^V\to F_3^V,\ F_2^V\to F_4^V)\,,\\
H^{(2)\,A}_{1/2,0} &= H^{(1)\,A}_{1/2,0}(F_1^A\to F_3^A,\ F_2^A\to F_4^A)\,,\\
H^{(1)\,V}_{1/2,t} &= -\frac{1}{\sqrt{q^2}}\Big\{\sqrt{Q_+}\Big[(m_{\Lambda_b}-m_\Lambda)F_1^V + q^2F_5^V\Big]\Big\}\,,\\
H^{(1)\,A}_{1/2,t} &= -\frac{1}{\sqrt{q^2}}\Big\{\sqrt{Q_-}\Big[(m_{\Lambda_b}+m_\Lambda)F_1^A - q^2F_5^A\Big]\Big\}\,,\\
H^{(2)\,V}_{1/2,t} &= H^{(1)\,V}_{1/2,t}(F_1^V\to F_3^V,\ F_5^V\to F_6^V)\,,\\
H^{(2)\,A}_{1/2,t} &= H^{(1)\,A}_{1/2,t}(F_1^A\to F_3^A,\ F_5^A\to F_6^A)\,,
\end{aligned}
\tag{2.17}
$$

where

$$
\begin{aligned}
Q_+ &= (m_{\Lambda_b}+m_\Lambda)^2 - q^2\,,\\
Q_- &= (m_{\Lambda_b}-m_\Lambda)^2 - q^2\,,
\end{aligned}
$$

and

$$
\begin{aligned}
F_1^V &= A_1 - D_1 + B_1 - E_1\,,\\
F_1^A &= A_1 - D_1 - B_1 + E_1\,,\\
F_2^V &= F_1^V(1\to 2)\,,\\
F_2^A &= F_1^A(1\to 2)\,,\\
F_3^V &= A_1 + D_1 + B_1 + E_1\,,\\
F_3^A &= A_1 + D_1 - B_1 - E_1\,,\\
F_4^V &= F_3^V(1\to 2)\,,\\
F_4^A &= F_3^A(1\to 2)\,,\\
F_5^V &= F_1^V(1\to 3)\,,\\
F_5^A &= F_1^A(1\to 3)\,,\\
F_6^V &= F_3^V(1\to 3)\,,\\
F_6^A &= F_3^A(1\to 3)\ .
\end{aligned}
\tag{2.18}
$$

The remaining helicity amplitudes can be obtained from the parity relations

$$
H^{V,(A)}_{-\lambda,-\lambda_W} = +(-)H^{V,(A)}_{\lambda,\lambda_W}\ . \tag{2.19}
$$

The square of the matrix element for the $\Lambda_b\to\Lambda\ell^+\ell^-$ decay is given as

$$
\begin{aligned}
|\mathcal{M}|^2 = &\left|\mathcal{M}^{++}_{+1/2}\right|^2 + \left|\mathcal{M}^{+-}_{+1/2}\right|^2 + \left|\mathcal{M}^{-+}_{+1/2}\right|^2 + \left|\mathcal{M}^{--}_{+1/2}\right|^2\\
&+ \left|\mathcal{M}^{++}_{-1/2}\right|^2 + \left|\mathcal{M}^{+-}_{-1/2}\right|^2 + \left|\mathcal{M}^{-+}_{-1/2}\right|^2 + \left|\mathcal{M}^{--}_{-1/2}\right|^2\ .
\end{aligned}
\tag{2.20}
$$

Following the standard methods used in literature (see the third reference in [28]), the normalized joint angular decay distribution for the two cascade decay

$$\Lambda_b^{1/2^+} \to \Lambda^{1/2^+}\Big(\to a(1/2^+)+b(0^-)\Big)+V(\to \ell^+\ell^-)\,,$$

can be written as

$$\begin{aligned}\frac{d\Gamma}{dq^2 d\cos\theta\, d\cos\theta_\Lambda} &= \left|\frac{G\alpha}{4\sqrt{2}\pi}V_{tb}V_{ts}^*\frac{1}{2}\right|^2 \frac{\sqrt{\lambda(m_{\Lambda_b}^2,m_\Lambda^2,q^2)}\sqrt{\lambda(m_\Lambda^2,m_a^2,m_b^2)}}{1024\pi^3 m_{\Lambda_b}^3 m_\Lambda^2}\\ &\times v\mathcal{B}(\Lambda_b\to a+b)\left|\mathcal{M}\right|^2\,, \end{aligned} \tag{2.21}$$

where the polar angle θ_Λ is the angle of the $a(1/2^+)$ momentum in the rest frame of the Λ baryon. Note that in this expression we perform integration over the azimuthal angle φ between the planes of the two decays $\Lambda \to a+b$ and $V \to \ell^+\ell^-$. Our final result for the differential decay width is

$$\begin{aligned}\frac{d\Gamma}{dq^2 d\cos\theta\, d\cos\theta_\Lambda} &= \left|\frac{G\alpha}{4\sqrt{2}\pi}V_{tb}V_{ts}^*\frac{1}{2}\right|^2 \frac{\sqrt{\lambda(m_{\Lambda_b}^2,m_\Lambda^2,q^2)}\sqrt{\lambda(m_\Lambda^2,m_a^2,m_b^2)}}{1024\pi^3 m_{\Lambda_b}^3 m_\Lambda^2} v\mathcal{B}(\Lambda\to a+b)\\ &\Bigg\{(1+\alpha_\Lambda\cos\theta_\Lambda)\Bigg[\Big(8m_\ell^2\sin^2\theta\left|A_{+1/2,+1}\right|^2+(1-\cos\theta)^2q^2\left|A_{+1/2,+1}-vB_{+1/2,+1}\right|^2\\ &+(1+\cos\theta)^2q^2\left|A_{+1/2,+1}+vB_{+1/2,+1}\right|^2\Big)+8m_\ell^2\cos^2\theta\left|A_{+1/2,0}\right|^2+8m_\ell^2\left|B_{+1/2,t}\right|^2\\ &+\sin^2\theta q^2\Big(2\left|A_{+1/2,0}\right|^2+2v^2\left|B_{+1/2,0}\right|^2\Big)\Bigg]\\ &+(1-\alpha_\Lambda\cos\theta_\Lambda)\Bigg[\Big(8m_\ell^2\sin^2\theta\left|A_{-1/2,-1}\right|^2+(1+\cos\theta)^2q^2\left|A_{-1/2,-1}-vB_{-1/2,-1}\right|^2\\ &+(1-\cos\theta)^2q^2\left|A_{-1/2,-1}+vB_{-1/2,-1}\right|^2\Big)+8m_\ell^2\cos^2\theta\left|A_{-1/2,0}\right|^2+8m_\ell^2\left|B_{-1/2,t}\right|^2\\ &+\sin^2\theta q^2\Big(2\left|A_{-1/2,0}\right|^2+2v^2\left|B_{-1/2,0}\right|^2\Big)\Bigg]\Bigg\}\,. \end{aligned} \tag{2.22}$$

In eq. (2.22) we induce the following definitions:

$$\begin{aligned} H^{(1)}_{\lambda_i,\lambda_W}+H^{(2)}_{\lambda_i,\lambda_W} &= A_{\lambda_i,\lambda_W}\\ H^{(1)}_{\lambda_i,\lambda_W}-H^{(2)}_{\lambda_i,\lambda_W} &= B_{\lambda_i,\lambda_W}\,. \end{aligned} \tag{2.23}$$

One can easily see that in addition to the variables that exist in eq. (2.22), there appears a new variable θ_Λ and integration of eq. (2.22) over it gives the differential decay width for the $\Lambda_b \to \Lambda\ell^+\ell^-$ decay.

It is well known that heavy quarks $b(c)$ resulting from Z decay are polarized. It is shown in [44, 45] that a sizeable fraction of the b quark polarization retained in fragmentation of heavy quarks to heavy baryons. Therefore, an additional set of polarization observables can be obtained if the polarization of the heavy Λ_b baryon is taken into account.

In order to take polarization of the Λ_b baryon into consideration we will use the density matrix method. The spin density matrix of Λ baryon is

$$\rho = \frac{1}{2}\begin{pmatrix} 1+\mathcal{P}\cos\theta_\Lambda^S & \mathcal{P}\sin\theta_\Lambda^S \\ \mathcal{P}\sin\theta_\Lambda^S & 1-\mathcal{P}\cos\theta_\Lambda^S \end{pmatrix}, \tag{2.24}$$

where $\mathcal{P}$ is the polarization of Λ_b, and θ_Λ^S is the angle that the polarization of Λ_b makes with the momentum of Λ, in the rest frame of Λ_b.

The four-fold decay distribution can easily be obtained from eq. (2.22). Obviously, there appears on the left-hand side of eq. (2.22) the distribution over θ_Λ^S, i.e., $d/d\cos\theta_\Lambda^S$. Hence the right-hand side of the same equation can be modified as follows:

$$\begin{aligned}
&|+1/2,+1|^2 \to (1-\mathcal{P}\cos\theta_\Lambda^S)\,|+1/2,+1|^2\,,\\
&\Big\{\,|+1/2,t|^2,\ |+1/2,0|^2,\ (+1/2,t)(+1/2,0)^*\Big\} \to (1+\mathcal{P}\cos\theta_\Lambda^S)\Big\{\,|+1/2,t|^2\,, |+1/2,0|^2\,,\\
&(+1/2,t)(+1/2,0)^*\Big\}\,,\\
&\Big\{(+1/2,+1)(+1/2,t)^*,\ (+1/2,+1)(+1/2,0)^*\Big\} \to \mathcal{P}\sin\theta_\Lambda^S\Big\{(+1/2,+1)(+1/2,t)^*,\\
&(+1/2,1)(+1/2,0)^*\Big\}\,,\\
&|-1/2,-1|^2 \to (1+\mathcal{P}\cos\theta_\Lambda^S)\,|-1/2,-1|^2\,,\\
&\Big\{\,|-1/2,t|^2,\ |-1/2,0|^2,\ (-1/2,t)(-1/2,0)^*\Big\} \to (1-\mathcal{P}\cos\theta_\Lambda^S)\Big\{\,|-1/2,t|^2,\ |-1/2,0|^2\,,\\
&(-1/2,t)(-1/2,0)^*\Big\}\,,\\
&\Big\{(-1/2,-1)(-1/2,t)^*,\ (-1/2,-1)(-1/2,0)^*\Big\} \to \mathcal{P}\sin\theta_\Lambda^S\Big\{(-1/2,-1)(-1/2,t)^*,\\
&(-1/2,-1)(-1/2,0)^*\Big\}\,.
\end{aligned} \tag{2.25}$$

It follows from eqs. (2.22) and (2.24) that the cascade decay $\Lambda_b \to \Lambda(\to a+b)\,V^*(\to \ell^+\ell^-)$ has a rich angular structure. Therefore, study of different distributions will prove useful in separating various angular coefficients in the experiments. For this reason, instead of analyzing the full four-fold angular distributions, one can investigate the individual angular distributions and their relations to the new Wilson coefficients. Foe example, the polar angle θ_Λ distribution of the cascade decay $\Lambda \to a+b$ can be obtained from eq. (2.22) by performing integration over θ, as a result of which takes the form

$$\frac{d\Gamma}{dq^2\,d\cos\theta_\Lambda} \sim 1+\alpha\alpha_\Lambda\cos\theta_\Lambda\,, \tag{2.26}$$

where the asymmetry parameter α is defined as

$$\begin{aligned}
\alpha = \frac{8}{3\Delta}\Big\{&4m_\ell^2\left|A_{+1/2,+1}\right|^2 + 2q^2\Big(\left|A_{+1/2,+1}\right|^2 + v^2\left|B_{+1/2,+1}\right|^2\Big)\\
&+2m_\ell^2\left|A_{+1/2,0}\right|^2 + 6m_\ell^2\left|B_{+1/2,t}\right|^2 + q^2\Big(\left|A_{+1/2,0}\right|^2 + v^2\left|B_{+1/2,0}\right|^2\Big)\\
&-4m_\ell^2\left|A_{-1/2,-1}\right|^2 - 2q^2\Big(\left|A_{-1/2,-1}\right|^2 + v^2\left|B_{-1/2,-1}\right|^2\Big)\\
&-q^2(\left|A_{-1/2,0}\right|^2 + v^2\left|B_{-1/2,0}\right|^2\Big) - 2m_\ell^2\left|A_{-1/2,0}\right|^2 - 6m_\ell^2\left|B_{-1/2,t}\right|^2\Big\}\,,
\end{aligned} \tag{2.27}$$

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where

$$\Delta = \frac{8}{3}\Big\{4m_\ell^2\left|A_{+1/2,+1}\right|^2 + 2q^2\Big(\left|A_{+1/2,+1}\right|^2 + v^2\left|B_{+1/2,+1}\right|^2\Big)$$
$$+2m_\ell^2\left|A_{+1/2,0}\right|^2 + 6m_\ell^2\left|B_{+1/2,t}\right|^2 + q^2\Big(\left|A_{+1/2,0}\right|^2 + v^2\left|B_{+1/2,0}\right|^2\Big)$$
$$+4m_\ell^2\left|A_{-1/2,-1}\right|^2 + 2q^2\Big(\left|A_{-1/2,-1}\right|^2 + v^2\left|B_{-1/2,-1}\right|^2\Big)$$
$$+q^2(\left|A_{-1/2,0}\right|^2 + v^2\left|B_{-1/2,0}\right|^2\Big) + 2m_\ell^2\left|A_{-1/2,0}\right|^2 + 6m_\ell^2\left|B_{-1/2,t}\right|^2\Big\}\,. \quad (2.28)$$

For the polar angle distribution in the cascade decay $V^* \to \ell^+\ell^-$ we integrate eq. (2.22) over $\theta_\Lambda$and we get

$$\frac{d\Gamma}{dq^2\,d\cos\theta} \sim 1 + 2\alpha_\theta\cos\theta + \beta_\theta\cos^2\theta\,, \quad (2.29)$$

where

$$\alpha_\theta = \frac{1}{\Delta_1}2vq^2\mathrm{Re}\Big[A_{+1/2,+1}B^*_{+1/2,+1} - A_{-1/2,-1}B^*_{-1/2,-1}\Big]\,, \quad (2.30)$$

$$\beta_\theta = \frac{1}{\Delta_1}\Big\{-4m_\ell^2\left|A_{+1/2,+1}\right|^2 + q^2\Big(\left|A_{+1/2,+1}\right|^2 + v^2\left|B_{+1/2,+1}\right|^2\Big)$$
$$+4m_\ell^2\left|A_{+1/2,0}\right|^2 - q^2\Big(\left|A_{+1/2,0}\right|^2 + v^2\left|B_{+1/2,0}\right|^2\Big)$$
$$-4m_\ell^2\left|A_{-1/2,-1}\right|^2 + q^2\Big(\left|A_{-1/2,-1}\right|^2 + v^2\left|B_{-1/2,-1}\right|^2\Big)$$
$$+4m_\ell^2\left|A_{-1/2,0}\right|^2 - q^2\Big(\left|A_{-1/2,0}\right|^2 + v^2\left|B_{-1/2,0}\right|^2\Big)\Big\}\,, \quad (2.31)$$

and

$$\Delta_1 = 4m_\ell^2\left|A_{+1/2,+1}\right|^2 + q^2\Big(\left|A_{+1/2,+1}\right|^2 + v^2\left|B_{+1/2,+1}\right|^2\Big)$$
$$+4m_\ell^2\left|B_{+1/2,t}\right|^2 + q^2\Big(\left|A_{+1/2,0}\right|^2 + v^2\left|B_{+1/2,0}\right|^2\Big)$$
$$+4m_\ell^2\left|A_{-1/2,-1}\right|^2 + q^2\Big(\left|A_{-1/2,-1}\right|^2 + v^2\left|B_{-1/2,-1}\right|^2\Big)$$
$$+4m_\ell^2\left|B_{-1/2,t}\right|^2 + q^2\Big(\left|A_{-1/2,0}\right|^2 + v^2\left|B_{-1/2,0}\right|^2\Big)\,, \quad (2.32)$$

If the polarization of the initial Λ_b is considered, a new symmetry parameter, which depends on θ_Λ^S appears. Performing integrations over θ_Λ and θ, we get

$$\frac{d\Gamma}{dq^2\,d\cos\theta_\Lambda^S} \sim 1 - \alpha_{\Lambda_S}\mathcal{P}\cos\theta_\Lambda^S\,, \quad (2.33)$$

where

$$\alpha_{\Lambda_S} = \frac{8}{3\Delta}\Big\{4m_\ell^2\left|A_{+1/2,+1}\right|^2 + 2q^2\Big(\left|A_{+1/2,+1}\right|^2 + v^2\left|B_{+1/2,+1}\right|^2\Big)$$
$$-2m_\ell^2\left|A_{+1/2,0}\right|^2 - q^2\Big(\left|A_{+1/2,0}\right|^2 + v^2\left|B_{+1/2,0}\right|^2\Big) - 6m_\ell^2\left|B_{+1/2,t}\right|^2$$
$$-4m_\ell^2\left|A_{-1/2,-1}\right|^2 - 2q^2\Big(\left|A_{-1/2,-1}\right|^2 + v^2\left|B_{-1/2,-1}\right|^2\Big)$$
$$+2m_\ell^2\left|A_{-1/2,0}\right|^2 + q^2\Big(\left|A_{-1/2,0}\right|^2 + v^2\left|B_{-1/2,0}\right|^2\Big) + 6m_\ell^2\left|B_{-1/2,t}\right|^2\Big\}\,, \quad (2.34)$$

where Δ is given in eq. (2.28).

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3. Numerical analysis

In this section we present our numerical results for the asymmetry parameters α_θ, α_{θ_Λ}, $\alpha_{\theta^S_\Lambda}$ and β. The values of the input parameters we use in our calculations are: $|V_{tb}V^*_{ts}| = 0.0385$, $m_\tau = 1.77\,\mathrm{GeV}$, $m_\mu = 0.106\,\mathrm{GeV}$. $m_b = 4.8\,\mathrm{GeV}$. For the Wilson coefficients we use their SM values which are given as: $C_7^{SM} = -0.313$, $C_9^{SM} = 4.344$ and $C_{10}^{SM} = -4.669$. In further numerical analysis, the values of the new Wilson coefficients which describe new physics beyond the SM are needed. The Wilson coefficients C_{BR} and C_{SL} are strictly constrained from $b \to s\gamma$ decay. The SM prediction on the branching ratio for the $b \to s\gamma$ decay coincide, practically, with experimental result and there seems to be no noticeable deviation between them. Therefore we can fix the values of C_{BR} and C_{SL} by substituting their SM values, i.e., $C_{BR} = -2m_b C_7^{\mathrm{eff}}$, $C_{SL} = -2m_s C_7^{\mathrm{eff}}$, where $C_7^{\mathrm{eff}} = -0.313$. Furthermore some of the Wilson coefficients describing vector interactions are restricted strongly by the present experimental data on branching ratios for the $B \to K\ell^+\ell^-$ and $B \to K^*\ell^+\ell^-$ decays [25, 26]. Using the experimental result on branching ratios for the above-mentioned decays, we obtain the following restrictions on C_{LL} and C_{RL}: $-2 \leq C_{LL} \leq 0$, and $0 \leq C_{RL} \leq 2.3$. The remaining Wilson are all varied in the region $-\left|C_{10}^{SM}\right| \leq C_X \leq +\left|C_{10}^{SM}\right|$. The upper bound on branching ratio of $B_s \to \mu^+\mu^-$ [46] suggests that this is the right order of magnitude for the vector interaction coefficients.

Few words about the Wilson coefficients C_9^{eff} are in order. Note that $\mathcal{M}(b \to s\ell^+\ell^-)$ for the $b \to s\ell^+\ell^-$ decay, although being a free quark decay amplitude, contains certain long distance effects from matrix elements of the four quark operators $\langle \ell^+\ell^- s\,|\mathcal{O}_i|\,b\rangle$ (explicit form of the operators $\mathcal{O}_1$–$\mathcal{O}_6$ can be found in [47, 48]) which are usually combined with the coefficient C_9 in an "effective" Wilson coefficient. For this reason, in exclusive decays one can define C_9^{eff} as

$$C_9^{\mathrm{eff}}(m_b, \hat{s}) = C_9(m_b)\left[1 + \frac{\alpha_s(\mu)}{\pi}\omega(\hat{s})\right] + Y_{SD}(m_b, \hat{s})\,, \tag{3.1}$$

where Y_{SD} corresponds to the above-mentioned four quark operator matrix elements, and $w(\hat{s})$ represents $\mathcal{O}(\alpha_f)$ corrections coming from one-gluon exchange in the matrix elements of the corresponding operator, whose explicit form can be found in [47]. The perturbative calculation leads to the following result for $Y_{SD}(\hat{s}, m_b)$:

$$\begin{aligned} Y_{SD}(m_b, \hat{s}) = {} & g\left(\hat{m}_c, \hat{s}\right) C^{(0)} - \frac{1}{2} g\left(1, \hat{s}\right)\left[4C_3 + 4C_4 + 3C_5 + C_6\right] \\ & - \frac{1}{2} g\left(0, \hat{s}\right)\left[C_3 + 3C_4\right] + \frac{2}{9}\left[3C_3 + C_4 + 3C_5 + C_6\right]\,, \end{aligned}$$

where

$$C^{(0)} = 3C_1 + C_2 + 3C_3 + C_4 + 3C_5 + C_6\,,$$

and the function $g(m_q, s)$ stands for the loops of quarks with mass m_q at the dilepton invariant mass s. This function develops absorptive parts for dilepton energies $s = 4m_q^2$:

$$g\left(\hat{m}_q, \hat{s}\right) = -\frac{8}{9}\ln \hat{m}_q + \frac{8}{27} + \frac{4}{9}y_q - \frac{2}{9}\left(2 + y_q\right)\sqrt{\left|1 - y_q\right|}$$

	$F(0)$	a_F	b_F
F_1	0.462	−0.0182	−0.000176
F_2	−0.077	−0.0685	0.00146

Table 1: Form factors for $\Lambda_b \to \Lambda\ell^+\ell^-$ decay in a three parameter fit.

$$\times\left[\Theta(1-y_q)\left(\ln\frac{1+\sqrt{1-y_q}}{1-\sqrt{1-y_q}}-i\pi\right)+\Theta(y_q-1)\,2\,\arctan\frac{1}{\sqrt{y_q-1}}\right],$$

where $\hat{m}_q = m_q/m_b$ and $y_q = 4\hat{m}_q^2/\hat{s}$

In addition to these perturbative contributions C_9 also receives long distance contributions coming from the production of $\bar{c}c$ resonances at intermediate states. Their contributions are represented by Y_{LD}, which has the form:

$$Y_{LD}(\hat{s}) = \frac{3}{\alpha^2}C^{(0)}\sum_{V_i=\psi(1s),\ldots,\psi(6s)}\frac{\pi\kappa_i\Gamma\left(V_i\to\ell^+\ell^-\right)M_{V_i}}{\left(M_{V_i}^2-\hat{s}m_b^2-iM_{V_i}\Gamma_{V_i}\right)}, \qquad (3.2)$$

where κ_i are the Fudge factors (see for example [7]). In regard to the absorptive parts that $C_9^{\rm eff}$ develops, no new fermions are introduced, and hence no new sources for the additional absorptive parts in the Wilson coefficients occur. For this reason we will assume that all new Wilson coefficients are real.

From the expressions of asymmetries it follows that the form factors are the main and the most important input parameters necessary in the numerical calculations. The calculation of the form factors of $\Lambda_b \to \Lambda$ transition does not exist at present. But, we can use the results from QCD sum rules in corporation with HQET [41, 49]. We noted earlier that, HQET allows us to establish relations among the form factors and reduces the number of independent form factors into two. In [41, 49], the q^2 dependence of these form factors are given as follows

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

The values of the parameters $F(0)$, a_F and b_F are given in table 1.

Note that the first analysis of the HQET structure of the $\Lambda_Q \to \Lambda_q$ transition is performed in [50] (see also [51]).

In order to have an idea about the sensitivity of our results to the specific parametrization of the two form factors predicted by the QCD sum rules in corporation with the HQET, we also have used another parametrization of the form factors based on the pole model and compared the results of both models. The dipole form of the form factors predicted by the pole model are given as:

$$F_{1,2}(E_\Lambda) = N_{1,2}\left(\frac{\Lambda_{\rm QCD}}{\Lambda_{\rm QCD}+E_\Lambda}\right)^2,$$

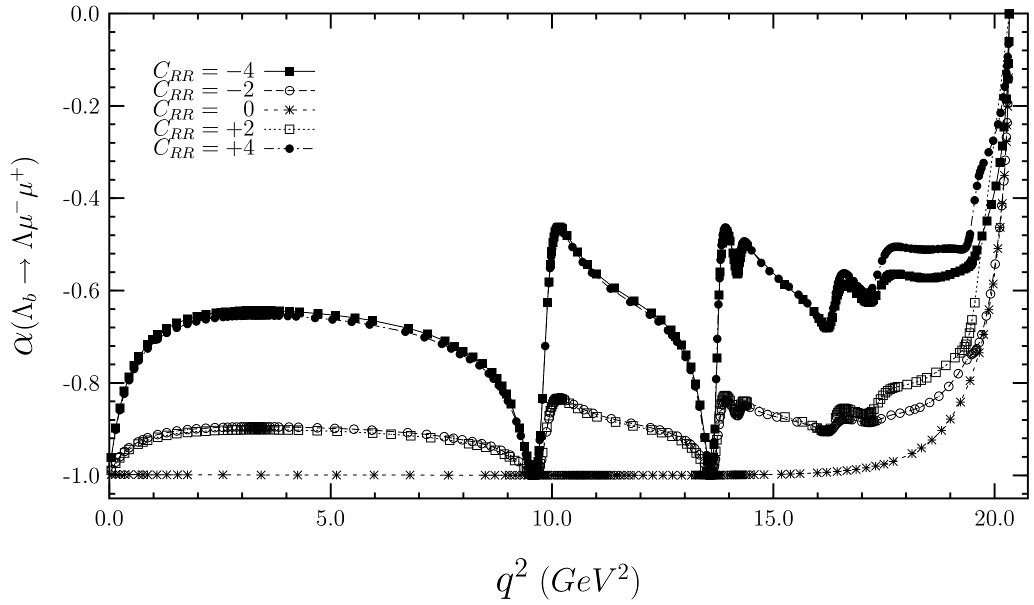

Figure 1: The dependence of the asymmetry parameter α on q^2 for the $\Lambda_b \to \Lambda\mu^+\mu^-$ decay, at five different fixed values of the vector type Wilson coefficient C_{RR}.

where

$$E_\Lambda = \frac{m_{\Lambda_b}^2 - m_\Lambda^2 - q^2}{2m_{\Lambda_b}},$$

and $\Lambda_{\rm QCD} = 0.2$, $|N_1| = 52.32$ and $|N_1| \simeq -0.25 N_1$ [52].

From the explicit expressions of the asymmetry parameters we see that they depend on the new Wilson coefficients and q^2. Therefore there might appear some difficulty in studying the dependence of the physical quantities on both variables in the experiments. For this reason we will study the dependence of the asymmetry parameters on q^2 at fixed values of the new Wilson coefficients.

In figure 1 we present the dependence of α on q^2 for the $\Lambda_b \to \Lambda\mu^+\mu^-$ decay at five fixed values of C_{RR}. We observe from this figure that at all values of C_{RR} the magnitude of α is smaller compared to the SM case for the whole range of q^2. The dependence of α on q^2 is not presented for the Wilson coefficients C_{LL} and C_{LR}, since our numerical analysis yields that α is not sensitive the presence of C_{LL} and C_{LR}, and it coincides with the SM result at all values of q^2.

In figure 2 we depict the dependence of α on q^2 for the $\Lambda_b \to \Lambda\mu^+\mu^-$ decay, at fixed values of C_{RL}. From this figure we see that, up to $q^2 = 18\,{\rm GeV}^2$ the magnitude of α is smaller compared to the SM prediction at $C_{RL} = 2$, but for $q^2 > 18\,{\rm GeV}^2$ the contribution of $C_{RL} = 2$ is exceeds that of the contribution of $C_{RL} = 0$ (i.e., SM case). In other words, investigation of α on q^2 in different kinematical regions of q^2 can give valuable information not only about the existence of the new physics, but also about the sign of the new Wilson coefficient C_{RL}.

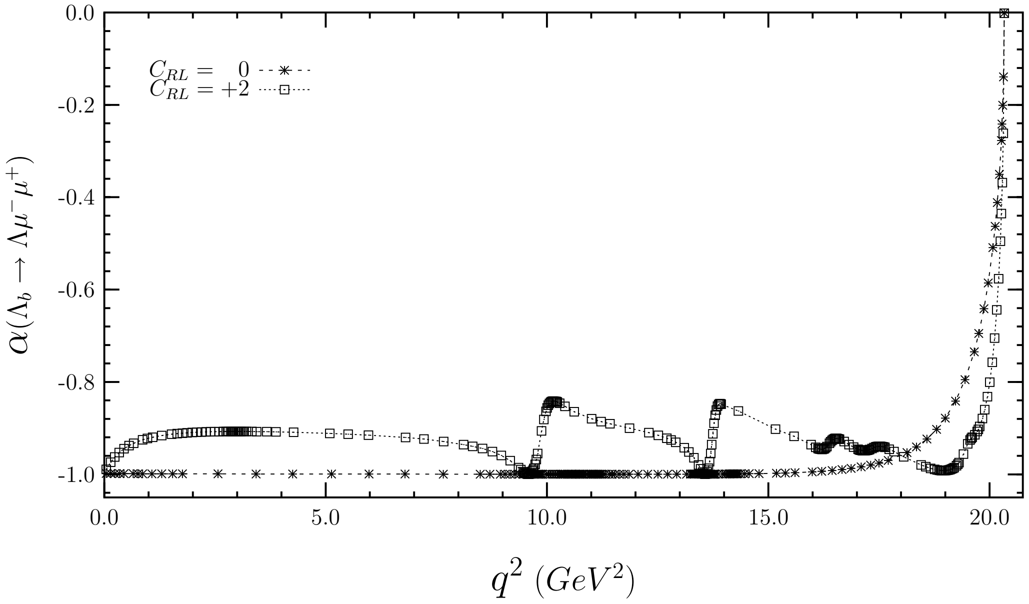

Figure 2: The same as in figure 1, but for the coefficient C_{RL}.

The study of the dependence of α on q^2 at fixed values of the new Wilson coefficients for the $\Lambda_b \to \Lambda \tau^+ \tau^-$ decay leads to the following results:

- The dependence of the asymmetry parameter α on q^2 is not sensitive to the presence of C_{LL} and C_{LR}, and practically there seems to be no departure from the SM prediction.

- The situation drastically changes in the presence of Wilson coefficients C_{RL} and C_{RR}. When C_{RR} $(C_{RL}) = 4(2)$, up to the range $q^2 = 18\,\mathrm{GeV}^2$, the value of α is two times smaller (as modulo) compared to that SM prediction; and when C_{RR} $(C_{RL}) = -4(2)$ the departure from from the SM result is about 50% (30%) . Therefore measurement of the asymmetry parameter α at different values of q^2 can give useful hint about the existence of C_{RL} and C_{RR}.

Next, we analyze the dependence of α_θ and β_θ for the $\Lambda_b \to \Lambda \mu^+ \mu^-$ decay. Our results can be summarized as follows:

- The zero of position of α_θ is shifted to the right (left) (see figures 3 and 4) when C_{LL} is negative (C_{LR} is positive). The essential point here is that, similar to the $B \to K^* \ell^+ \ell^-$ decay, the zero of position of α_θ is independent of the long distance effects and determined solely by short distance dynamics only.

- The zero of position of α_θ is practically independent of C_{RR} and C_{RL}.

Therefore, determination of zero position of α_θ can serve as a good tool for establishing the new physics beyond the SM, as well as the sign of new Wilson coefficients, which is controlled by the short distance physics only.

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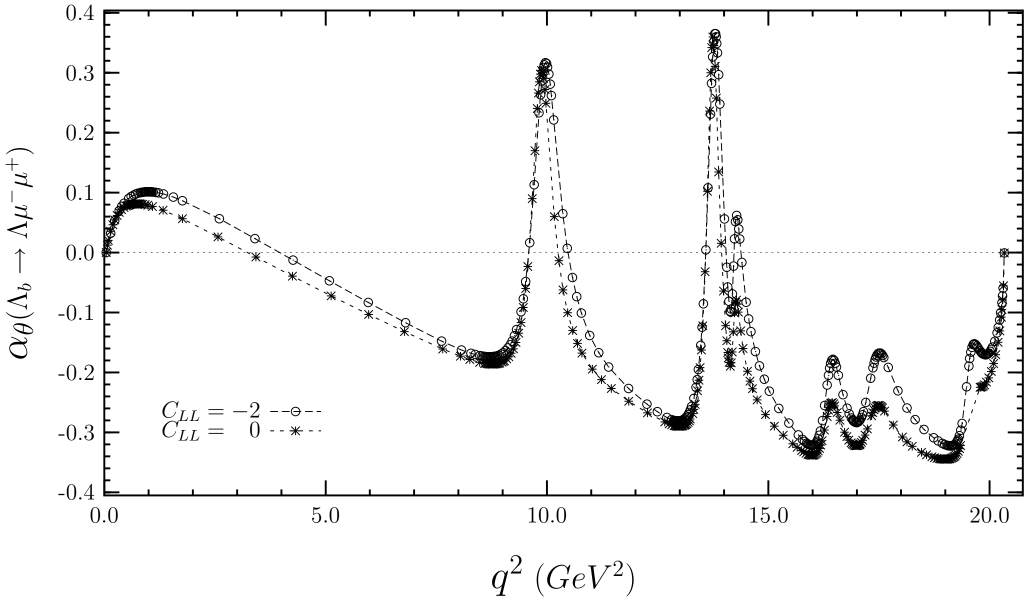

Figure 3: The dependence of the asymmetry parameter α_θ on q^2 for the $\Lambda_b \to \Lambda\mu^+\mu^-$ decay, at five different fixed values of the vector type Wilson coefficient C_{LL}.

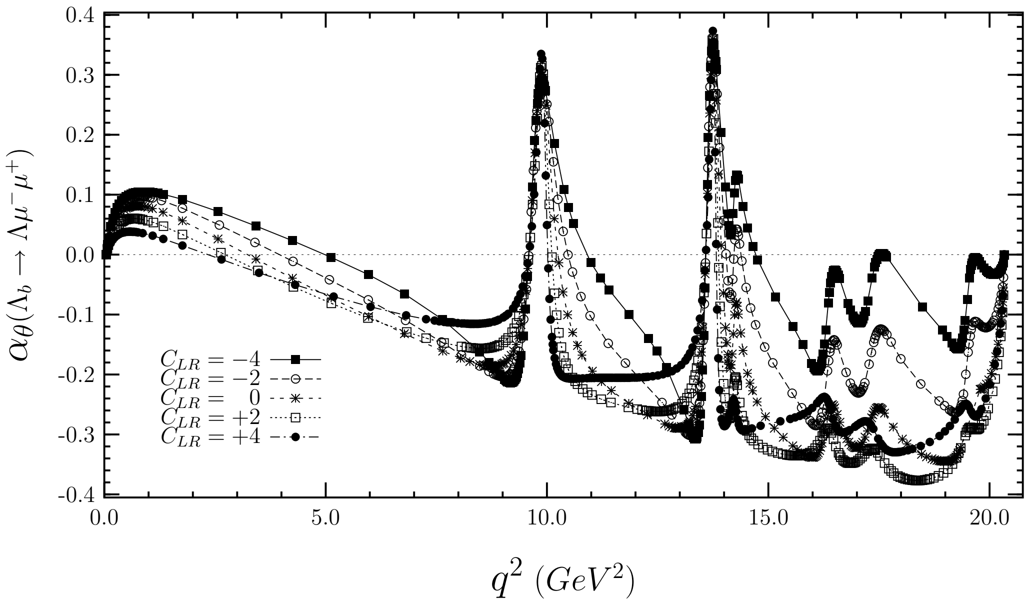

Figure 4: The same as in figure 3, but for the coefficient C_{LR}.

Moreover, the present analysis shows that β_θ is sensitive to the existence of the vector interactions C_{LL} in the region $1\,\mathrm{GeV}^2 \leq q^2 \leq 3\,\mathrm{GeV}^2$, and C_{LR} in the region $1\,\mathrm{GeV}^2 \leq q^2 \leq 8\,\mathrm{GeV}^2$. Therefore an investigation on the asymmetry parameter β_θ can give useful information about the existence of the vector interaction realized by the Wilson coefficients C_{LL} and C_{LR}. β_θ is not sensitive to the remaining two vector interactions C_{RL} and C_{RR} for the $\Lambda_b \to \Lambda\mu^+\mu^-$ decay.

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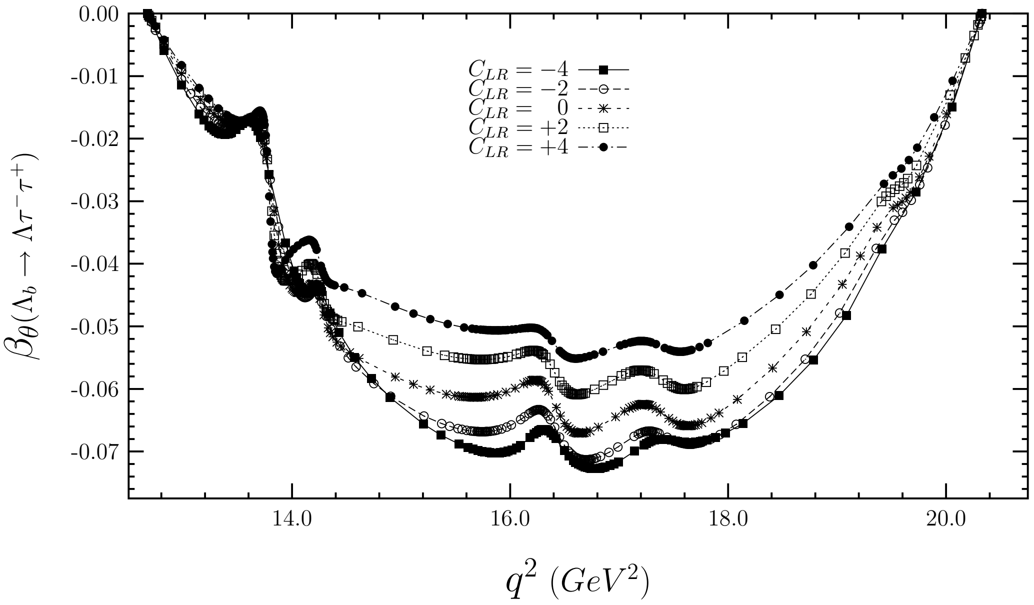

Figure 5: The dependence of the asymmetry parameter β_θ on q^2 for the $\Lambda_b \to \Lambda\tau^+\tau^-$ decay, at five different fixed values of the vector type Wilson coefficient C_{LR}.

From the analysis of the dependence of α_θ and β_θ on q^2 for the $\Lambda_b \to \Lambda\tau^+\tau^-$ decay we get:

- α_θ shows strong dependence on all Wilson coefficients.
- The dependence of β_θ on q^2 is similar to the SM case and at all values of all new Wilson coefficients the sign of β_θ is the same as in the SM case. Far from the resonance regions, it is strongly dependent on C_{LR}. For example, at $C_{LR} = \pm 4$, the departure from the SM result is about 50% larger when $14\,\mathrm{GeV}^2 \leq q^2 \leq 16\,\mathrm{GeV}^2$ (see figure 5).
- At positive (negative) values of C_{LR}, the magnitude of β_θ is smaller (larger) compared to that of SM prediction.

Finally, let us discuss the dependence of the asymmetry parameter α_{Λ_S} on q^2 at fixed values of C_X. In the $\Lambda_b \to \Lambda\mu^+\mu^-$ decay, α_{Λ_S} is more sensitive to all vector interactions in the kinematical region $1\,\mathrm{GeV}^2 \leq q^2 \leq 5\,\mathrm{GeV}^2$ (see figures 6 and 7, respectively).

For the $\Lambda_b \to \Lambda\tau^+\tau^-$ decay α_{Λ_S} is sensitive to all type of vector interactions (see figures 8–11), and it exhibits different behavior in its dependence on the new Wilson coefficients.

The dependence of α_{Λ_S} on the Wilson coefficients C_{LR} is similar to its dependence on C_{LL}. In the same region of q^2 when $C_{LR} = \pm 4$, α_{Λ_S} is two times smaller compared to that of the SM result. Note that near the end of the spectrum, i.e., $17.6\,\mathrm{GeV}^2 \leq q^2 \leq 19.6\,\mathrm{GeV}^2$, α_{Λ_S} changes its sign when $C_{LR} = +4$ (see figure 9). For all other type of vector interactions, the asymmetry parameter α_{Λ_S} does not change its sign.

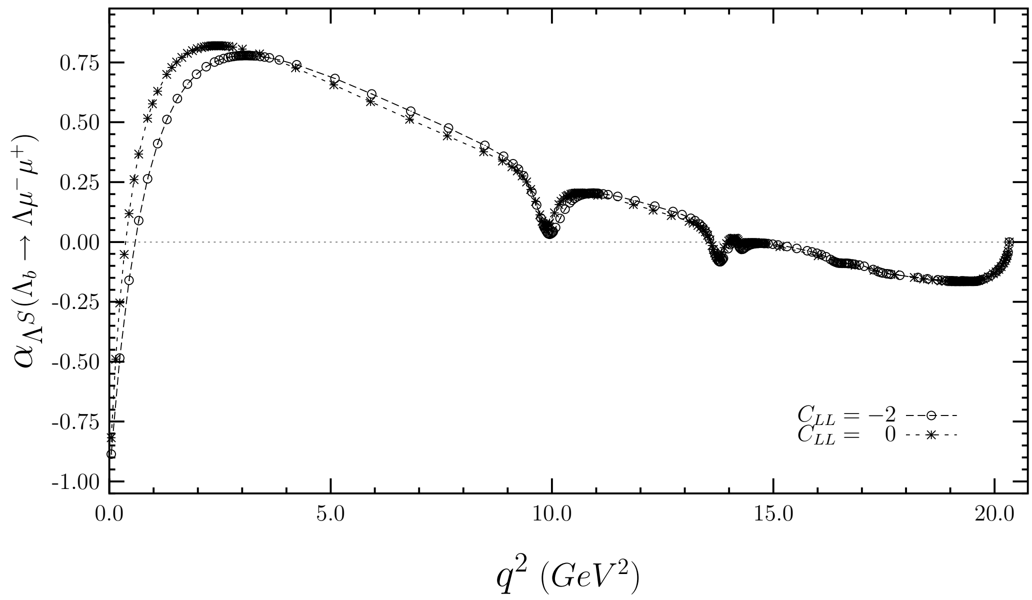

Figure 6: The dependence of the asymmetry parameter α_{Λ_S} on q^2 for the $\Lambda_b \to \Lambda\mu^+\mu^-$ decay, at five different fixed values of the vector type Wilson coefficient C_{LL}.

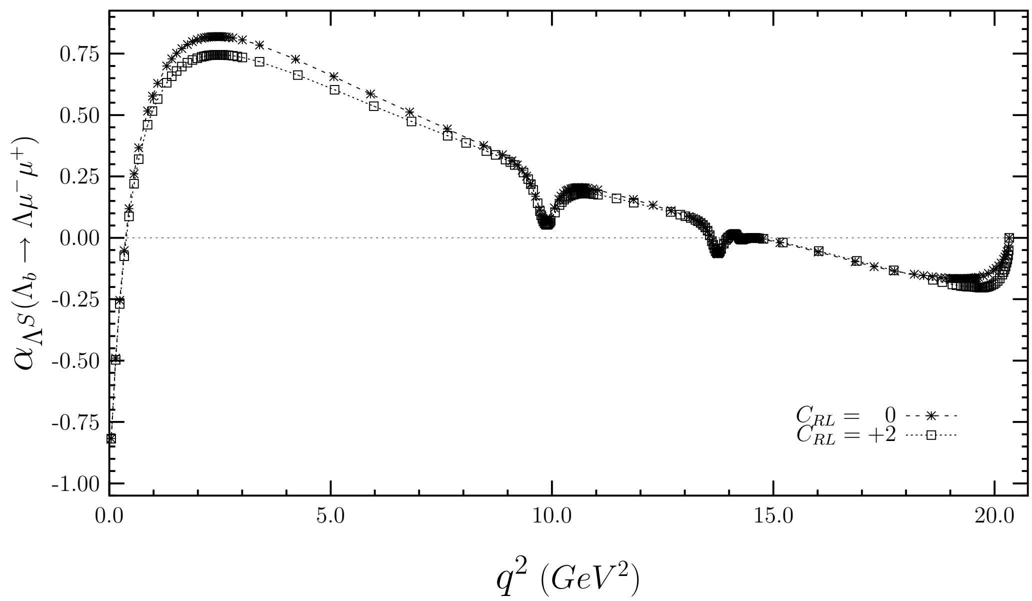

Figure 7: The same as in figure 6, but for the coefficient C_{RL}.

Therefore determination of the values of α_{Λ_S} in experiments can serve as an efficient tool for establishing the existence of the new type of vector interactions and also their signs.

In conclusion, in the present work we calculate the helicity amplitudes in the $\Lambda_b \to \Lambda\ell^+\ell^-$ decay in the framework of the minimal extension of the standard model with the inclusion of the new vector interactions. We analyze various asymmetry parameters of the

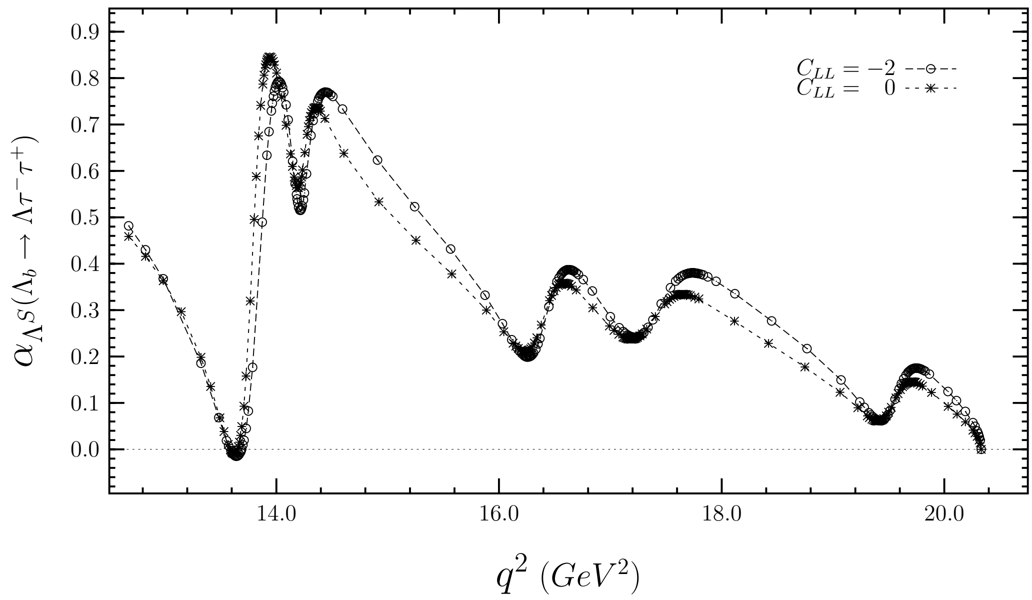

Figure 8: The same as in figure 6, but for $\Lambda_b \to \Lambda\tau^+\tau^-$ decay.

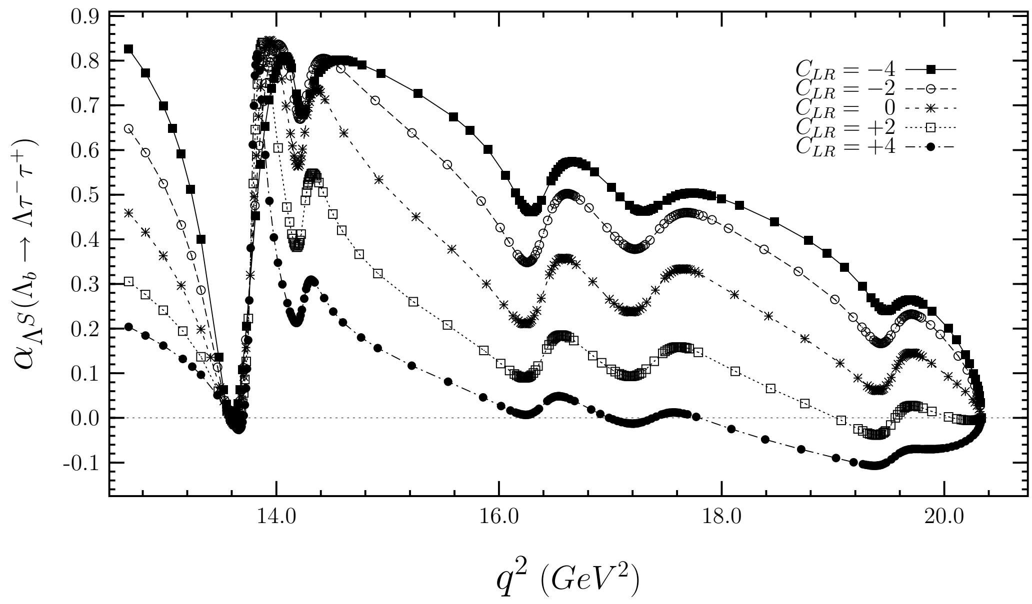

Figure 9: The same as in figure 8, but for the coefficient C_{LR}.

$\Lambda_b \to \Lambda(\to a + b)\, V^*(\to \ell^+\ell^-)$ decay with polarized and unpolarized heavy baryons and study their dependence on q^2 at fixed values of the new vector type interaction Wilson coefficients. We considered different asymmetry parameters and obtain that they exhibit strong dependence on different new Wilson coefficients. Therefore measurement of the different asymmetry parameters, namely, α, α_θ, β_θ and α_{Λ_S}, can give conformative informative about the existence of the new physics beyond the SM.

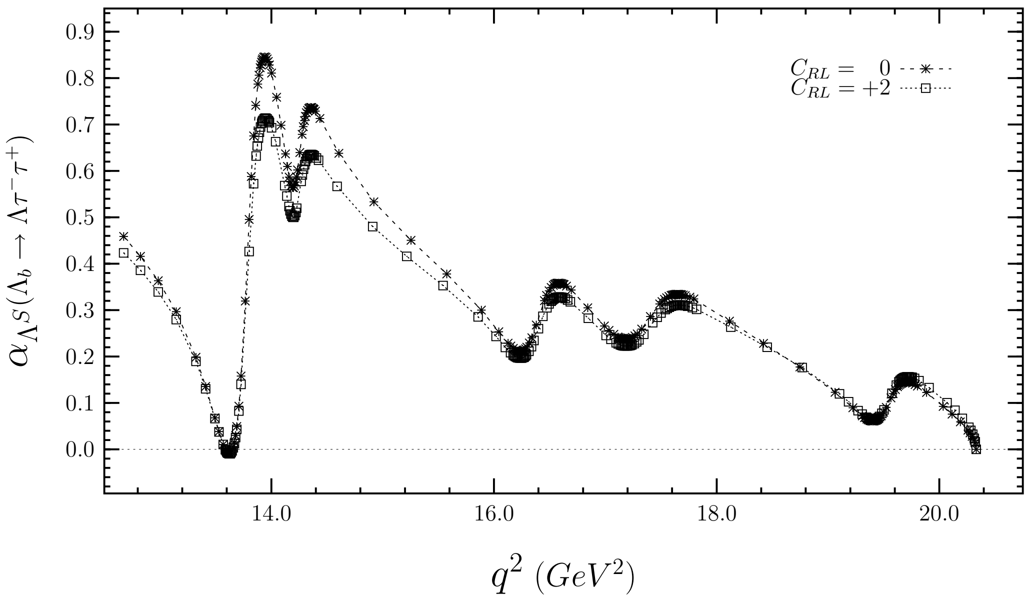

Figure 10: The same as in figure 8, but for the coefficient C_{RL}.

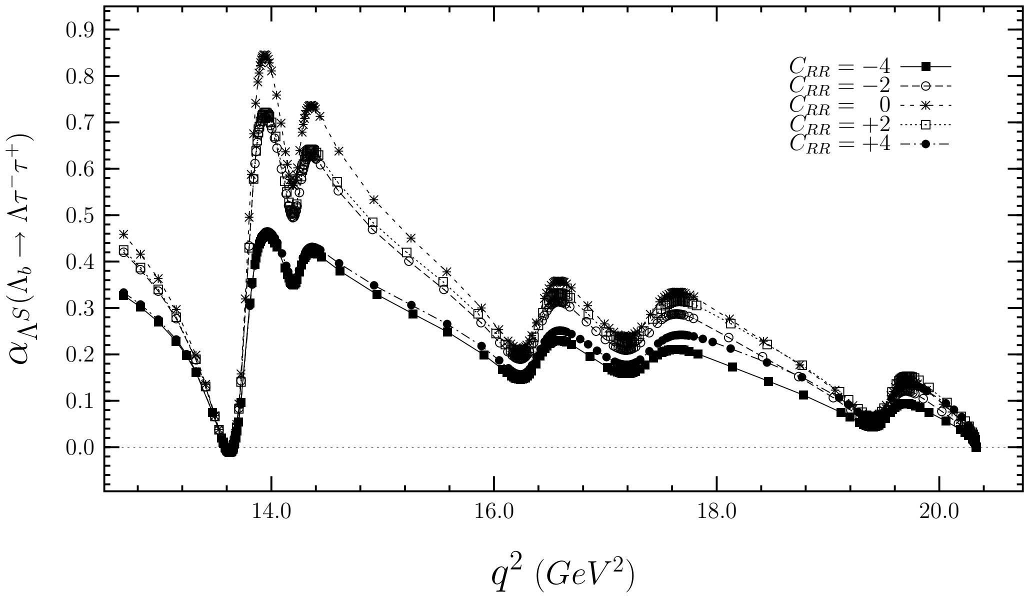

Figure 11: The same as in figure 8, but for the coefficient C_{RR}.

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