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Possible manifestation of new scalar interaction in P-odd asymmetries in $\Lambda_b \to \Lambda\tau^+\tau^-$ decay

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Abstract. Using the helicity amplitude method and including a new scalar type interaction in the matrix element of the exclusive semileptonic $\Lambda_\mathrm{b} \to \Lambda\tau^+\tau^-$ decay, P-odd asymmetries with polarized and unpolarized heavy baryons are investigated. The result is obtained that the study of P-odd asymmetries can be promising for establishing the new scalar sector beyond the SM.

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1 Introduction

Rare B decays induced by the flavor-changing neutral currents (FCNC) are promising candidates for testing the standard model (SM) at loop level and looking for new physics beyond the SM, such as the two Higgs doublet and supersymmetry. Moreover, these decays are also quite sensitive to the existence of new physics beyond the SM, since new particles running at loops can give contributions to these decays. New physics appears in rare decays through the Wilson coefficients which can take values different from their SM counterpart or through the new operator structures in the effective Hamiltonian (see for example [1] and references therein).

The exclusive decay which is described at the inclusive level by the $b \to s\ell^+\ell^-$ transition is the baryonic $\Lambda_b \to \Lambda\ell^+\ell^-$ decay. Unlike the mesonic $B \to K\ell^+\ell^-$, $B \to K^*\ell^+\ell^-$ decays, baryonic decays could maintain the helicity structure of the effective Hamiltonian for the $b \to s$ transition [2]. 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 studied in detail in the literature [3–8]. The experimental status of the heavy baryon decays is discussed in [9, 10].

In the present work we study different P-odd asymmetries in looking for new physics in the baryonic $\Lambda_b \to \Lambda\ell^+\ell^-$ decay due to the scalar interaction. In this analysis we use the helicity amplitude formalism and the polarization density matrix method (see the first and third references in [3]) to analyze the joint decay distributions in this decay. Note that the polarization effects due to the vector type interaction for the $\Lambda_b \to \Lambda\ell^+\ell^-$ decay are studied in the framework of SM and beyond in [1].

The paper is organized as follows. In Sect. 2, the matrix element for the $\Lambda_b \to \Lambda\tau^+\tau^-$ is calculated using the amplitude for the $b \to s$ transition that includes new scalar type interactions. In Sect. 3 the expressions for the different P-odd asymmetries are obtained.

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

In this section we derive the matrix element for the $\Lambda_b \to \Lambda\tau^+\tau^-$ decay which is described by the $b \to s\tau^+\tau^-$ transition. The amplitude for the $b \to s\tau^+\tau^-$ transition can be written in terms of the twelve model independent four-Fermi interactions as follows [11]:

$$\begin{aligned}
\mathcal{M} = \frac{G\alpha}{\sqrt{2}\pi} V_{tb}V_{ts}^* \Big\{ & C_{SL}\bar{s}_R \mathrm{i}\sigma_{\mu\nu}\frac{q^\nu}{q^2} b_L \bar{\tau}\gamma^\mu\tau \\
& + C_{BR}\bar{s}_L \mathrm{i}\sigma_{\mu\nu}\frac{q^\nu}{q^2} b_R \bar{\tau}\gamma^\mu\tau + C_{LL}^{\mathrm{tot}}\bar{s}_L\gamma_\mu b_L\bar{\tau}_L\gamma^\mu\tau_L \\
& + C_{LR}^{\mathrm{tot}}\bar{s}_L\gamma_\mu b_L\bar{\tau}_R\gamma^\mu\tau_R + C_{RL}\bar{s}_R\gamma_\mu b_R\bar{\tau}_L\gamma^\mu\tau_L \\
& + C_{RR}\bar{s}_R\gamma_\mu b_R\bar{\tau}_R\gamma^\mu\tau_R + C_{LRLR}\bar{s}_L b_R\bar{\tau}_L\tau_R \\
& + C_{RLLR}\bar{s}_R b_L\bar{\tau}_L\tau_R + C_{LRRL}\bar{s}_L b_R\bar{\tau}_R\tau_L \\
& + C_{RLRL}\bar{s}_R b_L\bar{\tau}_R\tau_L + C_T\bar{s}\sigma_{\mu\nu}b\bar{\tau}\sigma^{\mu\nu}\tau \\
& + \mathrm{i}C_{TE}\epsilon_{\mu\nu\alpha\beta}\bar{s}\sigma^{\mu\nu}b\bar{\tau}\sigma^{\alpha\beta}\tau \Big\}, \qquad (1)
\end{aligned}$$

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,

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which correspond to $-2m_s C_7^{\text{eff}}$ and $-2m_b C_7^{\text{eff}}$ in the SM, respectively. The $B \to K^* \gamma$ and $b \to s\gamma$ decays put stringent restrictions on C_{BR} and C_{SL}, and we will choose their SM values. The strange quark mass is neglected throughout in this work, i.e., we use $C_{BR} = -2m_b C_7^{\text{eff}}$ and $C_{SL} = 0$. The next four terms in (1) with coefficients C_{LL}^{tot}, C_{LR}^{tot}, C_{RL} and C_{RR} describe vector type interactions, two (C_{LL}^{tot} and C_{LR}^{tot}) of which contain SM contributions in the form $C_9^{\text{eff}} - C_{10}$ and $C_9^{\text{eff}} + C_{10}$, respectively. Thus, C_{LL}^{tot} and C_{LR}^{tot} can be written

$$\begin{aligned} C_{LL}^{\text{tot}} &= C_9^{\text{eff}} - C_{10} + C_{LL}\,, \\ C_{LR}^{\text{tot}} &= C_9^{\text{eff}} + C_{10} + C_{LR}\,, \end{aligned} \tag{2}$$

where C_{LL} and C_{LR} describe the contributions of new physics. Additionally, (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}). Note that in the present work we restrict ourselves to the minimal extension of the SM, i.e., we neglect the new vector and tensor type interactions. It should be noted here that theoretically our assumption can approximately be realized in the two-Higgs-doublet and supersymmetric models.

The amplitude of the exclusive $\Lambda_b \to \Lambda \tau^+ \tau^-$ decay is obtained by sandwiching (1) between the initial and final baryon states, $\langle \Lambda | \mathcal{M} | \Lambda_b \rangle$. It follows from (1) that the matrix elements

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

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

These matrix elements parametrized in terms of the form factors are as follows [12, 13]:

$$\langle \Lambda | \bar{s}\gamma_\mu b | \Lambda_b \rangle = \bar{u}_\Lambda \Big[f_1 \gamma_\mu + \mathrm{i} f_2 \sigma_{\mu\nu} q^\nu + f_3 q_\mu \Big] u_{\Lambda_b}\,, \tag{3}$$

$$\langle \Lambda | \bar{s}\gamma_\mu \gamma_5 b | \Lambda_b \rangle = \bar{u}_\Lambda \Big[g_1 \gamma_\mu \gamma_5 + \mathrm{i} g_2 \sigma_{\mu\nu} \gamma_5 q^\nu + g_3 q_\mu \gamma_5 \Big] u_{\Lambda_b}\,. \tag{4}$$

The form factors of the magnetic dipole operators are defined as

$$\begin{aligned} \langle \Lambda | \bar{s} \mathrm{i} \sigma_{\mu\nu} q^\nu b | \Lambda_b \rangle &= \bar{u}_\Lambda \Big[f_1^{\mathrm{T}} \gamma_\mu + \mathrm{i} f_2^{\mathrm{T}} \sigma_{\mu\nu} q^\nu + f_3^{\mathrm{T}} q_\mu \Big] u_{\Lambda_b}\,, \\ \langle \Lambda | \bar{s} \mathrm{i} \sigma_{\mu\nu} \gamma_5 q^\nu b | \Lambda_b \rangle &= \bar{u}_\Lambda \Big[g_1^{\mathrm{T}} \gamma_\mu \gamma_5 + \mathrm{i} g_2^{\mathrm{T}} \sigma_{\mu\nu} \gamma_5 q^\nu + g_3^{\mathrm{T}} q_\mu \gamma_5 \Big] \\ &\quad \times u_{\Lambda_b}\,. \end{aligned} \tag{5}$$

The matrix elements of scalar and pseudoscalar operators can be obtained by multiplying both sides of (3) and (4) with q_μ and using the equation of motion, as a result of which we get

$$\begin{aligned} \langle \Lambda | \bar{s} b | \Lambda_b \rangle &= \frac{1}{m_b - m_s} \bar{u}_\Lambda \Big[f_1 (m_{\Lambda_b} - m_\Lambda) + f_3 q^2 \Big] u_\Lambda\,, \\ \langle \Lambda | \bar{s} \gamma_5 b | \Lambda_b \rangle &= \frac{1}{m_b + m_s} \bar{u}_\Lambda \Big[g_1 (m_{\Lambda_b} + m_\Lambda) \gamma_5 - g_3 q^2 \Big] u_\Lambda\,. \end{aligned}$$

Using these definitions of the form factors, for the matrix element of $\Lambda_b \to \Lambda \tau^+ \tau^-$ we get [14, 15]

$$\begin{aligned} \mathcal{M} = \frac{G\alpha}{4\sqrt{2}\pi} V_{tb} V_{ts}^* \frac{1}{2} \Big\{ &\bar{\tau}\gamma_\mu (1-\gamma_5) \tau\, \bar{u}_\Lambda \\ &\times \Big[(A_1 - D_1) \gamma_\mu (1+\gamma_5) + (B_1 - E_1) \gamma_\mu (1-\gamma_5) \\ &+ \mathrm{i} \sigma_{\mu\nu} q^\nu \Big((A_2 - D_2)(1+\gamma_5) + (B_2 - E_2)(1-\gamma_5) \Big) \\ &+ q_\mu \Big((A_3 - D_3)(1+\gamma_5) + (B_3 - E_3)(1-\gamma_5) \Big) \Big] u_{\Lambda_b} \\ &+ \bar{\tau}\gamma_\mu (1+\gamma_5) \tau\, \bar{u}_\Lambda \\ &\times \Big[(A_1 + D_1) \gamma_\mu (1+\gamma_5) + (B_1 + E_1) \gamma_\mu (1-\gamma_5) \\ &+ \mathrm{i} \sigma_{\mu\nu} q^\nu \Big((A_2 + D_2)(1+\gamma_5) + (B_2 + E_2)(1-\gamma_5) \Big) \\ &+ q_\mu \Big((A_3 + D_3)(1+\gamma_5) + (B_3 - E_3)(1-\gamma_5) \Big) \Big] u_{\Lambda_b} \\ &+ \frac{1}{2} \bar{\tau} (1-\gamma_5) \tau\, \bar{u}_\Lambda \Big[(P_1 - P_2 + R_1 - R_1)(1-\gamma_5) \\ &+ (P_1 + P_2 - R_1 - R_2)(1+\gamma_5) \Big] u_{\Lambda_b} \\ &+ \frac{1}{2} \bar{\tau} (1+\gamma_5) \tau\, \bar{u}_\Lambda \Big[(P_1 - P_2 + R_2 - R_1)(1-\gamma_5) \\ &+ (P_1 + P_2 + R_1 + R_2)(1+\gamma_5) \Big] u_{\Lambda_b} \Big\}\,, \end{aligned} \tag{6}$$

where

$$\begin{aligned} A_1 &= \frac{1}{q^2} \left(f_1^{\mathrm{T}} + g_1^{\mathrm{T}} \right) (-2m_b C_7^{\text{eff}}) + (f_1 - g_1)\, C_9^{\text{eff}}\,, \\ A_2 &= A_1\,(1 \to 2)\,, \\ A_3 &= A_1\,(1 \to 3)\,, \\ B_1 &= A_1\,(g_1 \to -g_1;\; g_1^{\mathrm{T}} \to -g_1^{\mathrm{T}})\,, \\ B_2 &= B_1\,(1 \to 2)\,, \\ B_3 &= B_1\,(1 \to 3)\,, \\ D_1 &= C_{10}\,(f_1 - g_1)\,, \\ D_2 &= D_1\,(1 \to 2)\,, \\ D_3 &= D_1\,(1 \to 3)\,, \\ E_1 &= D_1\,(g_1 \to -g_1)\,, \\ E_2 &= E_1\,(1 \to 2)\,, \\ E_3 &= E_1\,(1 \to 3)\,, \\ P_1 &= \frac{1}{m_b} \Big(f_1\,(m_{\Lambda_b} - m_\Lambda) + f_3 q^2 \Big) \\ &\quad \times \Big(C_{LRLR} + C_{RLLR} + C_{LRRL} + C_{RLRL} \Big)\,, \\ P_2 &= P_1\,(C_{LRRL} \to -C_{LRRL};\; C_{RLRL} \to -C_{RLRL})\,, \\ R_1 &= \frac{1}{m_b} \Big(g_1\,(m_{\Lambda_b} + m_\Lambda) - g_3 q^2 \Big) \\ &\quad \times \Big(C_{LRLR} - C_{RLLR} + C_{LRRL} - C_{RLRL} \Big)\,, \\ R_2 &= R_1\,(C_{LRRL} \to -C_{LRRL};\; C_{RLRL} \to -C_{RLRL})\,. \end{aligned} \tag{7}$$

From (6) we see that the $\Lambda_b \to \Lambda \tau^+ \tau^-$ decay is described in terms of many form factors. It is shown in [16] that heavy quark effective theory (HQET) reduces the number of independent form factors to two (F_1 and F_2)

irrespective of the Dirac structure of the corresponding operators, i.e.,

$$\left\langle \Lambda(p_\Lambda) \left| \bar{s}\Gamma b \right| \Lambda(p_{\Lambda_b}) \right\rangle = \bar{u}_\Lambda \Big[F_1(q^2) + \not{v} F_2(q^2) \Big] \Gamma u_{\Lambda_b} , \tag{8}$$

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 (3)–(5) with (8), one can easily obtain the following relations among them (see also [13, 14]):

$$\begin{aligned}
g_1 &= f_1 = f_2^{\mathrm{T}} = g_2^{\mathrm{T}} = F_1 + \sqrt{\hat{r}_\Lambda} F_2 , \\
g_2 &= f_2 = g_3 = f_3 = \frac{F_2}{m_{\Lambda_b}} , \\
g_1^{\mathrm{T}} &= f_1^{\mathrm{T}} = \frac{F_2}{m_{\Lambda_b}} q^2 , \\
g_3^{\mathrm{T}} &= \frac{F_2}{m_{\Lambda_b}} \left(m_{\Lambda_b} + m_\Lambda \right) , \\
f_3^{\mathrm{T}} &= -\frac{F_2}{m_{\Lambda_b}} \left(m_{\Lambda_b} - m_\Lambda \right) ,
\end{aligned} \tag{9}$$

where $\hat{r}_\Lambda = m^2_\Lambda / m^2_{\Lambda_b}$. It should be noted here that the first analysis of the HQET structure of the $\Lambda_Q \to \Lambda_q$ transitions has been performed in [17, 18]

In order to obtain the helicity amplitudes for the $\Lambda_b \to \Lambda\tau^+\tau^-$ decay, it is convenient to consider this decay as a cascade decay $\Lambda_b \to \Lambda V^* \to \Lambda\tau^+\tau^-$, where V^* is the off-shell γ or Z bosons.

The matrix element of $\Lambda_b \to \Lambda\tau^+\tau^-$ decay can be written in the following form:

$$\mathcal{M}^{\lambda_\tau \bar{\lambda}_\tau}_{\lambda_i} = \sum_{\lambda_{V^*}} \eta_{\lambda_{V^*}} L^{\lambda_\tau \bar{\lambda}_\tau}_{\lambda_{V^*}} H^{\lambda_i}_{\lambda_{V^*}} ,$$

where

$$L^{\lambda_\tau \bar{\lambda}_\tau}_{\lambda_{V^*}} = \varepsilon^\mu_{V^*} \left\langle \tau^-(p_\tau, \lambda_\tau)\, \tau^+(p_\tau, \bar{\lambda}_\tau) \left| J^\tau_\mu \right| 0 \right\rangle , \tag{10}$$

$$H^{\lambda_i}_{\lambda_{V^*}} = (\varepsilon^\mu_{V^*})^* \left\langle \Lambda(p_\Lambda, \lambda_\Lambda) \left| J^i_\mu \right| \Lambda_b(p_{\Lambda_b}) \right\rangle , \tag{11}$$

where $\varepsilon^\mu_{V^*}$ is the polarization vector of the virtual intermediate vector boson, λ_τ and $\bar{\lambda}_\tau$ are the helicities of lepton and antilepton, respectively. The metric tensor can be expressed in terms of the polarization vector of the virtual vector particle $\varepsilon_V = \varepsilon(\lambda_V)$ as follows:

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

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 [19, 20]). The upper indices in (10) and (11) correspond to the helicities of the leptons, and the lower ones correspond to the helicity of the Λ baryon. Moreover, J^τ_μ and J^i_μ in (10) and (11) are the leptonic and hadronic currents, respectively.

In the calculations of the leptonic and baryonic matrix elements, usually two different frames are used. The leptonic amplitude $L^{\lambda_\tau \bar{\lambda}_\tau}_{\lambda_{V^*}}$ is calculated in the rest frame of the virtual vector boson with 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 the Λ_b baryon.

After standard calculations, we get for the helicity amplitudes

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

$$\begin{aligned}
\mathcal{M}^{--}_{-1/2} &= 2m_\tau \sin\theta \Big(H^{(1)}_{-1/2,-1} + H^{(2)}_{-1/2,-1} \Big) \\
&- 2m_\tau \cos\theta \Big(H^{(1)}_{-1/2,0} + H^{(2)}_{-1/2,0} \Big) \\
&+ 2m_\tau \Big(H^{(1)}_{-1/2,t} - H^{(2)}_{-1/2,t} \Big) \\
&+ \frac{1}{2}\sqrt{q^2}\Big[(1-v) J^{(1)}_{-1/2,0} - (1+v) J^{(2)}_{-1/2,0} \Big] \\
&+ \frac{1}{2}\sqrt{q^2}\Big[(1-v) J^{(1)}_{-1/2,t} - (1+v) J^{(2)}_{-1/2,t} \Big] , \qquad (12)
\end{aligned}$$

where the superscripts in $\mathcal{M}$ correspond to the helicities of leptons and antileptons, $v = \sqrt{1-4m_\tau^2/q^2}$ is the τ lepton velocity, θ is the angle of the positron in the rest frame of the intermediate boson with respect to its helicity axes, and

$$H^{(i)}_{\lambda,\lambda_W} = H^{(i)\,V}_{\lambda,\lambda_V} \pm H^{(i)\,A}_{\lambda,\lambda_V} , \qquad (13)$$

with λ and λ_V being the helicities of the leptons and intermediate vector bosons, respectively, $\lambda_V = t$ is the scalar (zero) helicity component of the initial V meson, and $i = 1,2$. The helicity amplitudes $H^{(i)\binom{V}{A}}_{\lambda,\lambda_W}$ are given as

$$\begin{aligned}
H^{(1)\binom{V}{A}}_{1/2,1} &= -\sqrt{Q_\mp}\left[F_1^{\binom{V}{A}} \mp (m_{\Lambda_b} \pm m_\Lambda) F_2^{\binom{V}{A}} \right] , \\
H^{(2)\binom{V}{A}}_{1/2,1} &= H^{(1)\binom{V}{A}}_{1/2,1} \\
&\times \left(F_1^{\binom{V}{A}} \to F_3^{\binom{V}{A}} ,\; F_2^{\binom{V}{A}} \to F_4^{\binom{V}{A}} \right) , \\
H^{(1)\binom{V}{A}}_{1/2,0} &= -\frac{1}{\sqrt{q^2}} \\
&\times \left\{ \sqrt{Q_\mp}\left[(m_{\Lambda_b} \pm m_\Lambda) F_1^{\binom{V}{A}} \mp q^2 F_2^{\binom{V}{A}} \right] \right\} , \\
H^{(2)\binom{V}{A}}_{1/2,0} &= H^{(1)\binom{V}{A}}_{1/2,0} \\
&\times \left(F_1^{\binom{V}{A}} \to F_3^{\binom{V}{A}} ,\; F_2^{\binom{V}{A}} \to F_4^{\binom{V}{A}} \right) , \\
H^{(1)\binom{V}{A}}_{1/2,t} &= -\frac{1}{\sqrt{q^2}} \\
&\times \left\{ \sqrt{Q_\pm}\left[(m_{\Lambda_b} \mp m_\Lambda) F_1^{\binom{V}{A}} \pm q^2 F_5^{\binom{V}{A}} \right] \right\} , \\
H^{(2)\binom{V}{A}}_{1/2,t} &= H^{(1)\binom{V}{A}}_{1/2,t} \left(F_1^{\binom{V}{A}} \to F_3^{\binom{V}{A}} ,\; F_5^{\binom{V}{A}} \to F_6^{\binom{V}{A}} \right) ,
\end{aligned}$$

$$\begin{aligned}
J^{(1)}_{+1/2,0} &= J^{(1)}_{+1/2,t} = \sqrt{Q_+}(P_1 - P_2) - \sqrt{Q_-}(R_1 - R_2) , \\
J^{(2)}_{+1/2,0} &= J^{(2)}_{+1/2,t} = J^{(1)}_{+1/2,0}(P_2 \to -P_2,\; R_2 \to -R_2) , \\
J^{(1)}_{-1/2,0} &= J^{(1)}_{-1/2,t} = J^{(1)}_{+1/2,0}(\sqrt{Q_-} \to -\sqrt{Q_-}) , \\
J^{(2)}_{-1/2,0} &= J^{(2)}_{-1/2,t} = J^{(2)}_{+1/2,0}(\sqrt{Q_-} \to -\sqrt{Q_-}) , \qquad (14)
\end{aligned}$$

where

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

and

$$\begin{aligned}
F_1^{V(A)} &= A_1 - D_1 + (-)(B_1 - E_1) , \\
F_2^{V(A)} &= F_1^{V(A)}(1 \to 2) , \\
F_3^{V(A)} &= A_1 + D_1 + (-)(B_1 + E_1) , \\
F_4^{V(A)} &= F_3^{V(A)}(1 \to 2) , \\
F_5^{V(A)} &= F_1^{V(A)}(1 \to 3) , \\
F_6^{V(A)} &= F_4^{V(A)}(2 \to 3) . \qquad (15)
\end{aligned}$$

The remaining helicity amplitudes can be obtained from the parity relations

$$H^{\binom{V}{A}}_{-\lambda,-\lambda_W} = \pm H^{\binom{V}{A}}_{\lambda,\lambda_W} . \qquad (16)$$

Following the standard methods used in the literature (see the third reference in [9]), the normalized joint angular decay distribution for the two cascade decay

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

is

$$\begin{aligned}
\frac{\mathrm{d}\Gamma}{\mathrm{d}q^2 \mathrm{d}\cos\theta\, \mathrm{d}\cos\theta_\Lambda} &= \left| \frac{G\alpha}{4\sqrt{2}\pi} V_{tb}V_{ts}^* \frac{1}{2} \right|^2 \\
&\times \frac{\sqrt{\lambda\left(m^2_{\Lambda_b}, m^2_\Lambda, q^2\right)}\sqrt{\lambda\left(m^2_\Lambda, m^2_a, m^2_b\right)}}{1024\pi^3 m^3_{\Lambda_b} m^2_\Lambda} v\mathcal{B}(\Lambda \to a + b) \\
&\times \Bigg\{ (1+\alpha_\Lambda \cos\theta_\Lambda)\Bigg[\Big(8m_\tau^2 \sin^2\theta \left|A_{+1/2,+1}\right|^2 \\
&+ (1-\cos\theta)^2 q^2 \left|A_{+1/2,+1} - vB_{+1/2,+1}\right|^2 \\
&+ (1+\cos\theta)^2 q^2 \left|A_{+1/2,+1} + vB_{+1/2,+1}\right|^2 \Big) \\
&+ 8m_\tau^2 \cos^2\theta \left|A_{+1/2,0}\right|^2 + 8m_\tau^2 \left|B_{+1/2,t}\right|^2 \\
&+ \sin^2\theta q^2 \Big(2\left|A_{+1/2,0}\right|^2 + 2v^2 \left|B_{+1/2,0}\right|^2 \Big) \\
&- 4m_\tau\sqrt{q^2}\Big(\mathrm{Re}\left[B_{+1/2,t}\left(D^*_{+1/2,t} + D^*_{+1/2,0}\right)\right] \\
&+ v\cos\theta \mathrm{Re}\left[A_{+1/2,0}\left(C^*_{+1/2,t} + C^*_{+1/2,0}\right)\right] \Big) \\
&+ \frac{q^2}{2}\Big(\left|D_{+1/2,t} + D_{+1/2,0}\right|^2 + v^2 \left|C_{+1/2,t} + C_{+1/2,0}\right|^2 \Big) \Bigg]
\end{aligned}$$

$$
\begin{aligned}
&+(1-\alpha_\Lambda\cos\theta_\Lambda)\Bigg[\Big(8m_\tau^2\sin^2\theta\left|A_{-1/2,-1}\right|^2\\
&+(1+\cos\theta)q^2\left|A_{-1/2,-1}-vB_{-1/2,-1}\right|^2\\
&+(1-\cos\theta)^2q^2\left|A_{-1/2,-1}+vB_{-1/2,-1}\right|^2\Big)\\
&+8m_\tau^2\cos^2\theta\left|A_{-1/2,0}\right|^2+8m_\tau^2\left|B_{-1/2,t}\right|^2\\
&+\sin^2\theta q^2\Big(2\left|A_{-1/2,0}\right|^2+2v^2\left|B_{-1/2,0}\right|^2\Big)\\
&-4m_\tau\sqrt{q^2}\Big(\mathrm{Re}\left[B_{-1/2,t}\left(D^*_{-1/2,t}+D^*_{-1/2,0}\right)\right]\\
&+v\cos\theta\mathrm{Re}\left[A_{-1/2,0}\left(C^*_{-1/2,t}+C^*_{-1/2,0}\right)\right]\Big)\\
&+\frac{q^2}{2}\Big(\left|D_{-1/2,t}+D_{-1/2,0}\right|^2\\
&+v^2\left|C_{-1/2,t}+C_{-1/2,0}\right|^2\Big)\Bigg]\Bigg\}\,. \qquad (17)
\end{aligned}
$$

In (17), the A_{λ_i,λ_W} are defined by

$$
\begin{aligned}
H^{(1)}_{\lambda_i,\lambda_W}+H^{(2)}_{\lambda_i,\lambda_W}&=A_{\lambda_i,\lambda_W}\,,\\
H^{(1)}_{\lambda_i,\lambda_W}-H^{(2)}_{\lambda_i,\lambda_W}&=B_{\lambda_i,\lambda_W}\,,\\
J^{(1)}_{\lambda_i,\lambda_W}+J^{(2)}_{\lambda_i,\lambda_W}&=C_{\lambda_i,\lambda_W}\,,\\
J^{(1)}_{\lambda_i,\lambda_W}-J^{(2)}_{\lambda_i,\lambda_W}&=D_{\lambda_i,\lambda_W}\,. \qquad (18)
\end{aligned}
$$

Note that in deriving (18), we perform an integration over the azimuthal angle φ between the planes of the two decays $\Lambda \to a+b$ and $V \to \tau^+\tau^-$.

It is well known that heavy quarks $b(c)$ resulting from Z decay are polarized. It is shown in [21, 22] that a sizeable fraction of the b quark polarization is retained in the 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 the polarization of the Λ_b baryon into consideration we will use the density matrix method. The spin density matrix of Λ baryon is

$$
\rho=\frac{1}{2}\begin{pmatrix}1+\mathcal{P}\cos\theta_\Lambda^S & \mathcal{P}\sin\theta_\Lambda^S\\ \mathcal{P}\sin\theta_\Lambda^S & 1-\mathcal{P}\cos\theta_\Lambda^S\end{pmatrix}, \qquad (19)
$$

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 (17), and when we use the density matrix, the right-hand side of this equation is modified:

$$
\begin{aligned}
&|1/2,1|^2 \to (1-\mathcal{P}\cos\theta_\Lambda^S)\,|1/2,1|^2\,,\\
&|1/2,t|^2\,,\ |1/2,0|^2\,,\ (1/2,0)(1/2,t)^*\\
&\quad\to(1+\mathcal{P}\cos\theta_\Lambda^S)\Big\{\,|1/2,0|^2\,,|1/2,t|^2\,,(1/2,0)(1/2,t)^*\Big\}\,,\\
&(1/2,1)(1/2,t)^*,\ (1/2,1)(1/2,0)^*\\
&\quad\to\mathcal{P}\sin\theta_\Lambda^S\Big\{(1/2,1)(1/2,0)^*,\ (1/2,1)(1/2,t)^*\Big\}\,,\\
&|-1/2,-1|^2\to(1+\mathcal{P}\cos\theta_\Lambda^S)\,|-1/2,-1|^2\,,\\
&|-1/2,t|^2\,,\ |-1/2,0|^2\,,\ (-1/2,0)(-1/2,t)^*\\
&\to(1-\mathcal{P}\cos\theta_\Lambda^S)\Big\{\,|-1/2,0|^2\,,\ |-1/2,t|^2\,,\\
&\qquad\qquad(-1/2,0)(-1/2,t)^*\Big\}\,,\\
&(-1/2,-1)(-1/2,-t)^*,\ (-1/2,-1)(-1/2,0)^*\\
&\to\mathcal{P}\sin\theta_\Lambda^S\times\Big\{(-1/2,-1)(-1/2,t)^*,\\
&\qquad\qquad(-1/2,-1)(-1/2,0)^*\Big\}\,. \qquad (20)
\end{aligned}
$$

From the expressions for the four-fold angular distribution we may define the following forward–backward asymmetries:

$$
\begin{aligned}
\mathcal{A}_\theta^{\mathrm{FB}}=&\left(\left[\int_0^{+1}\mathrm{d}\cos\theta-\int_{-1}^{0}\mathrm{d}\cos\theta\right]\int_{-1}^{+1}\mathrm{d}\cos\theta_\Lambda\right.\\
&\left.\times\int_{-1}^{+1}\mathrm{d}\cos\theta_\Lambda^S\frac{\mathrm{d}\Gamma}{\mathrm{d}q^2\,\mathrm{d}\cos\theta\,\mathrm{d}\cos\theta_\Lambda\,\mathrm{d}\cos\theta_\Lambda^S}\right)\\
&\Bigg/\left(\int_{-1}^{+1}\mathrm{d}\cos\theta\int_{-1}^{+1}\mathrm{d}\cos\theta_\Lambda\right.\\
&\left.\times\int_{-1}^{+1}\mathrm{d}\cos\theta_\Lambda^S\frac{\mathrm{d}\Gamma}{\mathrm{d}q^2\,\mathrm{d}\cos\theta\,\mathrm{d}\cos\theta_\Lambda\,\mathrm{d}\cos\theta_\Lambda^S}\right), \qquad (21)\\
\mathcal{A}_{\theta_\Lambda^S}^{\mathrm{FB}}=&\left(\left[\int_0^{+1}\mathrm{d}\cos\theta_\Lambda^S-\int_{-1}^{0}\mathrm{d}\cos\theta_\Lambda^S\right]\int_{-1}^{+1}\mathrm{d}\cos\theta\right.\\
&\left.\times\int_{-1}^{+1}\mathrm{d}\cos\theta_\Lambda\frac{\mathrm{d}\Gamma}{\mathrm{d}q^2\,\mathrm{d}\cos\theta\,\mathrm{d}\cos\theta_\Lambda\,\mathrm{d}\cos\theta_\Lambda^S}\right)\\
&\Bigg/\left(\int_{-1}^{+1}\mathrm{d}\cos\theta\int_{-1}^{+1}\mathrm{d}\cos\theta_\Lambda\right.\\
&\left.\times\int_{-1}^{+1}\mathrm{d}\cos\theta_\Lambda^S\frac{\mathrm{d}\Gamma}{\mathrm{d}q^2\,\mathrm{d}\cos\theta\,\mathrm{d}\cos\theta_\Lambda\,\mathrm{d}\cos\theta_\Lambda^S}\right), \qquad (22)\\
\mathcal{A}_{\theta_\Lambda}^{\mathrm{FB}}=&\left(\left[\int_0^{+1}\mathrm{d}\cos\theta_\Lambda-\int_{-1}^{0}\mathrm{d}\cos\theta_\Lambda\right]\int_{-1}^{+1}\mathrm{d}\cos\theta\right.\\
&\left.\times\int_{-1}^{+1}\mathrm{d}\cos\theta_\Lambda^S\frac{\mathrm{d}\Gamma}{\mathrm{d}q^2\,\mathrm{d}\cos\theta\,\mathrm{d}\cos\theta_\Lambda\,\mathrm{d}\cos\theta_\Lambda^S}\right)\\
&\Bigg/\left(\int_{-1}^{+1}\mathrm{d}\cos\theta\int_{-1}^{+1}\mathrm{d}\cos\theta_\Lambda\right.\\
&\left.\times\int_{-1}^{+1}\mathrm{d}\cos\theta_\Lambda^S\frac{\mathrm{d}\Gamma}{\mathrm{d}q^2\,\mathrm{d}\cos\theta\,\mathrm{d}\cos\theta_\Lambda\,\mathrm{d}\cos\theta_\Lambda^S}\right). \qquad (23)
\end{aligned}
$$

Using (21)–(23) one can easily find the explicit expessions for $\mathcal{A}_\theta^{\mathrm{AF}}$, $\mathcal{A}_{\theta_\Lambda^S}^{\mathrm{AF}}$ and $\mathcal{A}_{\theta_\Lambda}^{\mathrm{AF}}$.

3 Numerical analysis

In this section we will study the sensitivity of the P-odd asymmetries on the new scalar Wilson coefficients. The values of the input parameters we use in our calculations are $|V_{tb}V_{ts}^*| = 0.0385$, $m_\tau = 1.77\,\mathrm{GeV}$, $m_b = 4.8\,\mathrm{GeV}$, and we neglect the mass of the strange quark. In further numerical analysis, the values of the new Wilson coefficients which describe new physics beyond the SM are needed. In numerical calculations we will vary all new Wilson coefficients in the range $-\left|C_{10}^{\mathrm{SM}}\right| \le C_X \le \left|C_{10}^{\mathrm{SM}}\right|$. The experimental results on the branching ratio of the $B \to K^*\mu^+\mu^-$ decay [23] and the bound on the branching ratio of $B \to \mu^+\mu^-$ [24] suggest that this is the right order of magnitude for the vector and scalar interaction coefficients. For the values of the Wilson coefficients in the SM we use $C_7^{\mathrm{SM}} = -0.313$, $C_9^{\mathrm{SM}} = 4.344$ and $C_{10}^{\mathrm{SM}} = -4.669$. As far as the Wilson coefficient C_9^{SM} is considered, we take into account the short and the long distance contributions coming from the production of a $\bar{c}c$ pair at intermediate states (more about this issue can be found in [1]). It is well known 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 the $\Lambda_b \to \Lambda$ transition does not exist at present. But we can use the results from the QCD sum rules in co-operation with HQET [16, 25]. We noted earlier that HQET allows us to establish relations among the form factors and reduces the number of independent form factors to two. In [16, 25], the q^2 dependence of these form factors is 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.

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 difficulties in studying the dependence of the physical quantities on both variables in the experiments. In order to avoid this difficulty we perform our analysis at fixed values of C_X.

We see that $\mathcal{A}_\theta^{\mathrm{FB}}$ for the $\Lambda_b \to \Lambda\tau^+\tau^-$ decay is very sensitive to the existence of the scalar interaction with the coefficient C_{LRRL}, while it is independent of all other scalar interactions (see Fig. 1). Therefore, measurement of $\mathcal{A}_\theta^{\mathrm{FB}}(\Lambda_b \to \Lambda\tau^+\tau^-)$ can be quite informative for establishing the new scalar interactions.

Our numerical analysis shows that $\mathcal{A}_{\theta_\Lambda}^{\mathrm{FB}}$ for the $\Lambda_b \to \Lambda\tau^+\tau^-$ decay is rather sensitive to the scalar interactions with the coefficients C_{RLRL} and C_{RLLR}, while it is independent of the remaining scalar interaction coefficients. Close to the end of the allowed region ($q^2 > 18\,\mathrm{GeV}^2$), $\mathcal{A}_{\theta_\Lambda}^{\mathrm{FB}}(\Lambda_b \to \Lambda\tau^+\tau^-)$ shows a considerable departure from the SM result (see Fig. 2).

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

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

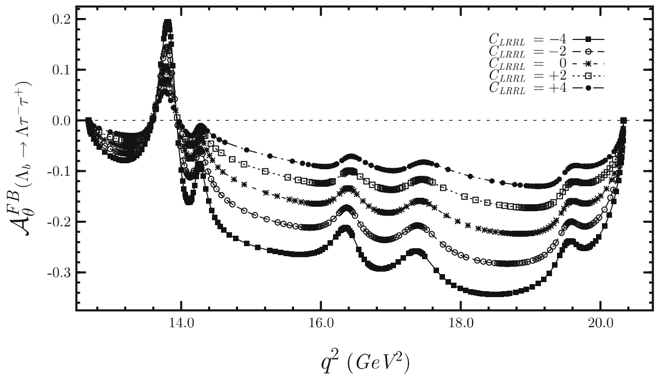

Fig. 1. The dependence of the P-odd forward–backward asymmetry $\mathcal{A}_\theta^{\mathrm{FB}}$ on q^2 at five different fixed values of the scalar type Wilson coefficient C_{LRRL} for the $\Lambda_b \to \Lambda\tau^+\tau^-$ decay

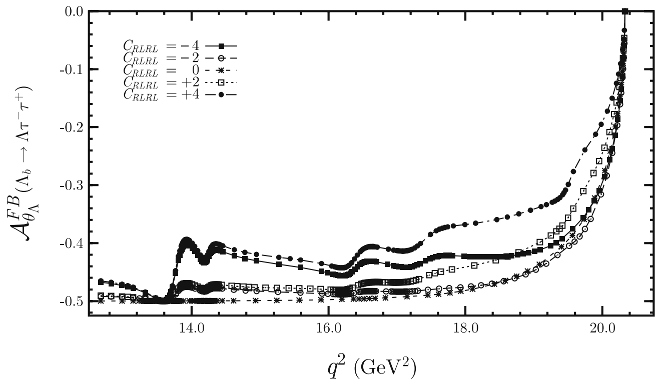

Fig. 2. The dependence of the P-odd forward–backward asymmetry $\mathcal{A}_{\theta_\Lambda}^{\mathrm{FB}}$ on q^2 at five different fixed values of the scalar type Wilson coefficient C_{RLRL} for the $\Lambda_b \to \Lambda\tau^+\tau^-$ decay

It is further observed that the $\mathcal{A}_{\theta_\Lambda^S}^{\mathrm{FB}}$ asymmetry for the $\Lambda_b \to \Lambda\tau^+\tau^-$ decay is very sensitive to the presence of the new scalar type interactions C_{LRRL} and C_{LRLR} (see Fig. 3). Since the dependence of $\mathcal{A}_{\theta_\Lambda^S}^{\mathrm{FB}}(\Lambda_b \to \Lambda\tau^+\tau^-)$ on the above-mentioned scalar coefficients turned out to be very similar, we present the one for the C_{LRRL} case.

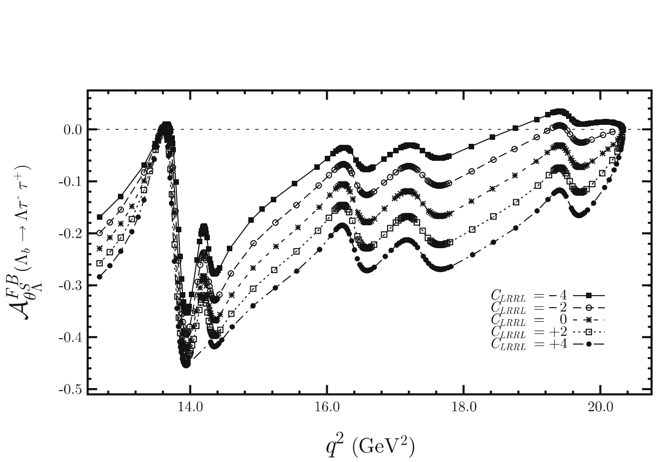

Fig. 3. The same as in Fig. 1, but for $\mathcal{A}_{\theta_\Lambda^S}^{\mathrm{FB}}$

From this figure we observe that there appears a new zero-position of $\mathcal{A}_{\theta_\Lambda^S}^{\mathrm{FB}}$, which is absent in the SM, and far from the resonance region, a considerable departure from the SM is predicted. For this reason study of the zero-position of $\mathcal{A}_{\theta_\Lambda^S}^{\mathrm{FB}}$ can give comprehensive information on the existence of new physics beyond the SM.

We can get additional information by measuring the magnitude of $\mathcal{A}_{\theta_\Lambda^S}^{\mathrm{FB}}$ in the region $18\,\mathrm{GeV}^2 \le q^2 \le 19.2\,\mathrm{GeV}^2$, which can be useful in determining the magnitude of the new Wilson coefficients.

In conclusion, we study the dependence of three P-odd forward–backward asymmetries on q^2 in the presence of the new scalar type interactions. The results we obtain can briefly be summarized as follows.

– Determination of the value of $\mathcal{A}_\theta^{\mathrm{FB}}$ for the $\Lambda_b \to \Lambda \tau^+ \tau^-$ decay can give invaluable information about the new physics, which is more sensitive to the presence of the scalar coefficient C_{LRRL}.
– It is shown that the P-odd asymmetry $\mathcal{A}_{\theta_\Lambda}^{\mathrm{FB}}$ for the $\Lambda_b \to \Lambda \tau^+ \tau^-$ decay exhibits a strong dependence on the scalar coefficients C_{RLRL} and C_{RLLR} as well.
– Our analysis predicts that, except for the resonance regions, there are new zero-points of $\mathcal{A}_{\theta_\Lambda^S}^{\mathrm{FB}}$ for the scalar coefficients C_{LRRL}, C_{LRLR} for the $\Lambda_b \to \Lambda \tau^+ \tau^-$ decay, which are absent in the SM. Therefore the determination of the zero-position of $\mathcal{A}_{\theta_\Lambda^S}$ can be very useful for establishing new physics beyond the SM due to the scalar interactions.

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