The $\Lambda_b \to \Lambda^*(1520)(\to N\bar{K})\ell^+\ell^-$ decay at low-recoil in HQET

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Abstract: In this paper we discuss the Standard Model and new physics sensitivity of the $\Lambda_b \to \Lambda^*(1520)(\to N\bar{K})\ell^+\ell^-$ decay at low-recoil, where $\ell^\pm$ are massive leptons and $N\bar{K} = \{pK^-, n\bar{K}^0\}$. We provide a full angular distribution with an operator basis that includes the Standard Model operator and its chirality flipped counterparts, and new scalar and pseudo-scalar operators. The resulting angular distribution allows us to construct observables that we study in the Standard Model and in model-independent new physics scenarios. To reduce the hadronic effects coming from the $\Lambda_b \to \Lambda^*(1520)$ transition form factors, we exploit the Heavy Quark Effective theory framework valid at low $\Lambda^*(1520)$ recoil, i.e., large dilepton invariant mass squared $q^2 \sim O(m_b^2)$. Working to the leading order in $1/m_b$ and including $O(\alpha_s)$ corrections, we compute the ‘improved Isgur-Wise relations’ between the form factors. The relations correlate the form factors and thereby allows the description of this decay at the low-recoil region with a smaller number of independent form factors.

Keywords: Rare Decays, Baryon Decays, Heavy Quark Effective Theory, New Physics
1 Introduction

A large number of experimental measurements in the $b \to s \ell^+ \ell^-$ transitions have shown deviations from the Standard Model (SM) predictions. The most recent of these deviations are the ratio of $B \to K^{(*)} \ell^+ \ell^-$ branching ratios for muon over electrons in the measurements of $R_K$ [1] and $R_{K^*}$ [2] which could be hints of lepton flavor universality (LFU) violation. The other deviations include the branching ratios of $B \to K\mu^+\mu^-$ [3], $B_s \to K^*\mu^+\mu^-$ [4, 5], $B_s \to \phi\mu^+\mu^-$ [6], and the optimized observables in the $B \to K^{(*)}\mu^+\mu^-$ decay [7]. These deviations have put the exploration of new physics (NP) in the $b \to s \ell^+ \ell^-$ transitions at the forefront of the $B$-physics program.

Indeed, the LHCb’s recent capability to study baryonic decays with acceptable statistic has opened up a new avenue to explore the $b \to s \ell^+ \ell^-$ transitions. For instance, the $\Lambda_b \to \Lambda(1116)\to N\pi\mu^+\mu^-$ has been measured by the LHCb [8, 9], and the observed branching ratio is lower than the SM predictions [10–13] with a trend that is also observed in the $B_s \to K^{(*)}\phi\mu^+\mu^-$ decays. The LHCb has also performed the LFU violation measurement in $\Lambda_b \to pK^-\ell^+\ell^-$ decay [14].

Theoretical studies of $\Lambda_b$ decays through a $b \to s \ell^+ \ell^-$ transition have received considerable attention in the recent years [10–13, 15–18]. Among the different semileptonic modes of $\Lambda_b$ decays to hadrons, the decay to $\Lambda^* \equiv \Lambda^*(1520)$ has the dominant contribution [19]. The $\Lambda^*$ has spin parity of $J^P = 3/2^-$ and decays strongly to $N\bar{K}$ pair. These characteristics make the $\Lambda^*$ easily distinguishable from the closely lying $\Lambda(1600)$, $\Lambda(1405)$, and the weakly decaying $\Lambda(1116)$, all of which have spin parity $J^P = 1/2^\pm$.

In this paper we discuss some aspects of $\Lambda_b \to \Lambda^*(\to N\bar{K})\ell^+\ell^-$ decay in the SM and beyond. The four-body decay proceeds through the weak decay of $\Lambda_b \to \Lambda^*\ell^+\ell^-$ and the subsequent strong decay $\Lambda^* \to N\bar{K}$. The SM operator basis for the weak decay is supplemented with its chirality flipped counterparts (henceforth called SM$'$ basis), and a scalar and pseudo-scalar operator basis (henceforth SP basis). We provide a full angular distribution and express the angular coefficients in terms of transversity amplitudes. The masses of the leptons in the final state are retained which can be important if the dilepton pair is heavy. From the four-fold angular distribution we construct several observables that we studied in the SM and in model-independent NP.

Like any exclusive decay, this mode also suffers from long-distance QCD dynamics coming from different sources. This include the “naively” factorizable contributions of the $b \to s \ell^+ \ell^-$ and $b \to s\gamma$ operators parametrized in terms of form factors. In two opposite kinematical regions, significant reduction to the number of independent form factors can be achieved by means of effective theories– the Soft-Collinear Effective Theory (SCET) [20–23] at large hadronic recoil or low dilepton invariant mass squared $q^2$, and the Heavy Quark Effective Theory (HQET) [24–26] at low-recoil or large $q^2$. A systematic discussions of the decay in the SCET framework [27–29] is beyond the scope of the present paper due to poor knowledge of the baryonic wave function and the fact that the spectator scattering effects are more complicated to deal with [30]. We will rather focus on the low-recoil region where the decay can be described by the HQET framework, and an operator product (OPE) expansion in $1/Q$ [31], where $Q \sim (m_b, \sqrt{q^2})$. The low-recoil OPE allows us to
include the contributions from the charm quark loops into the Wilson coefficients \[31\]. The HQET spin flavor symmetry ensures relations between the form factors, known as the Isgur-Wise relations \[32\]. Following the prescription of Ref. \[31\], we obtain ‘improved Isgur-Wise relations’ to leading order in \(1/m_b\) and including \(O(\alpha_s)\) corrections. By means of these relations, the form factors are correlated and each of the transversity amplitudes depend on a single form factor, as a result of which the short- and long-distance physics factorize in the angular coefficients. The reduction in the number of independent form factors lead to improved prediction in the low-recoil region. The low-recoil factorization allows us to construct observables that are sensitive to NP and have reduced dependence on form factor inputs.

The article is organized as follows. In the Sec. 2 we describe the effective Hamiltonian for the \(b \to s \ell^+ \ell^-\) transition in the SM+SM’+SP operator basis, and derive the transversity amplitudes. The four-fold angular distribution is worked out and the angular coefficients are derived in the Sec. 3. In Sec. 4 we construct the observables. In Sec. 5 we derive the improved Isgur-Wise relations between the the \(\Lambda_b \to \Lambda^*\) form factors, describe the low-recoil factorization, and construct observables for clean extraction of the electroweak physics. We present our numerical analysis in Sec. 6 and the summary is given in Sec. 7. We have also presented the necessary formulas in the appendices.

2 The Framework

2.1 Effective Hamiltonian

The \(\Lambda_b \to \Lambda^* \ell^+ \ell^-\) decay is governed by the \(b \to s \ell^+ \ell^-\) transition for which we assume the following effective Hamiltonian

\[ H_{\text{eff}} = - \frac{4G_F}{\sqrt{2}} V_{tb} V_{ts}^* \frac{\alpha_e}{4\pi} \left( \sum_i C_i O_i + \sum_j C_{j'} O_{j'}' \right), \]

where \(i = 7, 9, 10, S, P\) and \(j' = 7', 9', 10', S', P'.\) The operators read

\[ O_7 = \frac{m_b}{e} \left[ \bar{s} \sigma^{\mu\nu} P_R b \right] F_{\mu\nu}, \quad O_{7'} = \frac{m_b}{e} \left[ \bar{s} \sigma^{\mu\nu} P_L b \right] F_{\mu\nu}, \]

\[ O_9 = \left[ \bar{s} \gamma^\mu P_L b \right] \left[ \ell \gamma_\mu \ell' \right], \quad O_{9'} = \left[ \bar{s} \gamma^\mu P_R b \right] \left[ \ell \gamma_\mu \ell' \right], \]

\[ O_{10} = \left[ \bar{s} \gamma^\mu P_L b \right] \left[ \ell \gamma_\mu \gamma_5 \ell' \right], \quad O_{10'} = \left[ \bar{s} \gamma^\mu P_R b \right] \left[ \ell \gamma_\mu \gamma_5 \ell' \right], \]

\[ O_{S(\ell)} = \left[ \bar{s} P_R (L) b \right] \left[ \ell \ell' \right], \quad O_{P(\ell)} = \left[ \bar{s} P_R (L) b \right] \left[ \ell \gamma_5 \ell' \right]. \]

The operators \(O_{7,9,10}\) are the ones that appear in the SM and the rests appear in beyond the SM scenarios. NP operators of the the types \(O_{9,10}\) may also appear but the effects can be included trivially by the substitutions in the SM Wilson coefficients \(C_{9,10} \to C_{9,10} + \delta C_{9,10}^{NP}\). The Wilson coefficients \(C_{7,9}\) (usually denoted by \(C_{7,9}^{\text{eff}}\) in the literature) absorb the contributions of factorizable quark-loop contributions of the hadronic operators \(O_{1-6}\) and \(O_8\). For simplicity we have ignored the non-factorizable contributions of the operators \(O_{1-6}\) and \(O_8\) that are expected to play significant role in the large-recoil or low-\(q^2\) region \[22, 33\]. The rest of the parameters are as follows: \(G_F\) is the Fermi-constant, \(V_{tb}V_{ts}^*\) are the
Cabibbo-Kobayashi-Maskawa (CKM) elements, $\alpha_e = e^2/4\pi$ is the fine structure constant, and $P_{L(R)} = (1 + \gamma_5)/2$ are the chirality projectors. The $b$-quark mass appearing in the $O_{7,7'}$ operators are taken as the running mass in the modified minimal subtraction scheme ($\overline{\text{MS}}$). The contribution proportional to $V_{ub}V_{us}^*$ has been neglected since $V_{ub}V_{us}^* << V_{tb}V_{ts}^*$ and therefore, CP violation is absent in our analysis.

2.2 Decay kinematics

We assign the following momenta and spin variables to the different particles in the decay process

$$\Lambda_b(p, s_{\Lambda_b}) \to \Lambda^*(k, s_{\Lambda^*})\ell^+(q_1)\ell^-(q_2),$$

$$\Lambda^*(k, s_{\Lambda^*}) \to N(k_1, s_N)\bar{K}(k_2),$$

(i.e., $p, k, k_1, k_2, q_1$ and $q_2$ are the momenta of $\Lambda_b$, $\Lambda^*$, $N$, $\bar{K}$, and the positively and negatively charged leptons, respectively, and $s_{\Lambda_b, \Lambda^*, N}$ are the projections of the baryon spins on to the $z$-axis in their respective rest frames. For future convenience we define the momentum for dilepton pair $q^\mu = q_1^\mu + q_2^\mu$, (2.4)

and the momentum conservation gives $k^\mu = k_1^\mu + k_2^\mu$, $p^\mu = k^\mu + q^\mu$. The angle $\theta_\ell$ is defined as the one made by the lepton $\ell^-$ with respect to the $+z$ axis in the $\ell^+\ell^-$ rest frame, $\theta_{\Lambda^*}$ is the angle made by the nucleon with the $+z$ axis in the $N\bar{K}$ rest frame, and $\phi$ is the angle between the decay planes of the dielpton pair and the hadron pair. We have spelled out the kinematics in Appendix A.

2.3 The $\Lambda_b \to \Lambda^*\ell^+\ell^-$ decay

Assuming factorizations between the hadronic and the leptonic parts, the matrix element of the four-body decay $\Lambda_b \to \Lambda^* (\to N\bar{K})\ell^+\ell^-$ can be written as

$$\mathcal{M}(s_{\Lambda_b}, s_N, \lambda_1, \lambda_2) = \sum_{s_{\Lambda^*}} \mathcal{M}_{\Lambda_b}^{\lambda_1, \lambda_2}(s_{\Lambda_b}, s_{\Lambda^*}) \mathcal{M}_{\Lambda^*}(s_{\Lambda^*}, s_N),$$

(2.5)

where $\mathcal{M}_{\Lambda^*}(s_{\Lambda^*}, s_N)$ correspond to the matrix element for $\Lambda^* \to N\bar{K}$ which is discussed in the next section. The amplitudes for the first stage of the decay can be written as

$$\mathcal{M}_{\Lambda_b}^{\lambda_1, \lambda_2}(s_{\Lambda_b}, s_{\Lambda^*}) = \frac{4G_F}{\sqrt{2}} V_{tb} V_{ts}^* \frac{\alpha_e}{4\pi} \sum_{L(R)} \frac{1}{4} \left[ \sum_{\lambda} H_{VA,\lambda}^{L(R)}(s_{\Lambda_b}, s_{\Lambda^*}) L_{L(R),\lambda}^{\lambda_1, \lambda_2} \eta_{\lambda, \lambda} + H_{SP}^{L(R)}(s_{\Lambda_b}, s_{\Lambda^*}) L_{L(R)}^{\lambda_1, \lambda_2} \right],$$

(2.6)

where the hadronic and the leptonic helicity amplitudes are defined as the projection of the full hadronic and leptonic amplitudes on to the polarization direction of virtual gauge boson that decays to the dilepton pairs. Denoting the polarizations of the virtual gauge
boson as $\tilde{\epsilon}_\mu(\lambda)$ for different polarization states $\lambda = 0, \pm 1, t$, the leptonic helicity amplitudes are written as

$$L_{L(R)}^{\lambda_1, \lambda_2} = \langle \bar{\ell}(\lambda_1) \ell(\lambda_2) | \bar{\ell}(1 \mp \gamma_5) \ell | 0 \rangle,$$  
(2.7)

$$L_{L(R)}^{\lambda_1, \lambda_2} = \bar{\epsilon}^\mu(\lambda) \langle \bar{\ell}(\lambda_1) \ell(\lambda_2) | \bar{\ell} \gamma_\mu (1 \mp \gamma_5) \ell | 0 \rangle.$$ 
(2.8)

Our choices of the gauge boson polarizations and the expressions of $L_{L(R)}^{\lambda_1, \lambda_2}$ are detailed in the Appendix B and C, respectively.

The hadronic helicity amplitudes are similarly defined as

$$H_{VA,\lambda}^{L(R)}(s_{\Lambda b}, s_{\Lambda'}) = \bar{\epsilon}_s^\mu(\lambda) \langle \Lambda(k, s_{\Lambda'}) | \left( C_9 \mp C_{10} \right) \bar{s} \gamma^\mu (1 - \gamma_5) b + \left( C_{9'} \mp C_{10'} \right) \bar{s} \gamma^\mu (1 + \gamma_5) b$$

$$- \frac{2m_b}{q^2} \left( C_7 \bar{s} i q_6 \sigma^{\mu\nu}(1 + \gamma_5) b + C_7' \bar{s} i q_6 \sigma^{\mu\nu}(1 - \gamma_5) b \right) \right] \Lambda_b(p, s_{\Lambda_b}) \rangle,$$

(2.9)

$$H_{SP}^{L(R)}(s_{\Lambda b}, s_{\Lambda'}) = \langle \Lambda(k, s_{\Lambda'}) | \left( C_{9'} \mp C_{10'} \right) \bar{s}(1 - \gamma_5) b + \left( C_{9} \mp C_{10} \right) \bar{s}(1 + \gamma_5) b \right] \Lambda_b(p, s_{\Lambda_b}) \rangle.$$ 
(2.10)

In order to calculate the amplitudes, we need to know the form factor parametrization of the $\Lambda_b \to \Lambda^*$ hadronic matrix elements. We follow the helicity parametrization [17] (see Appendix D) where the $\langle \Lambda^* | \bar{s} \gamma^\mu | \Lambda_b \rangle$ ($\langle \Lambda^* | \bar{s} \gamma^\mu \gamma_5 | \Lambda_b \rangle$) transition is parametrized in terms of four $q^2$-dependent form factors $f_V^T$, $f_0^T$, $f_0^A$, $f_2^A$, $f_2^T$). The matrix elements of scalar and the pseudo-scalar currents $\langle \Lambda^* | \bar{s} b | \Lambda_b \rangle$, $\langle \Lambda^* | \bar{s} \gamma_5 b | \Lambda_b \rangle$ are obtained from the vector and axial vector matrix elements by the application of equation of motion and depend on $f_V^T$ and $f_2^T$ respectively. For the tensor currents the $\langle \Lambda^* | \bar{s} \sigma^{\mu\nu} q_6 b | \Lambda_b \rangle$ ($\langle \Lambda^* | \bar{s} \gamma_5 \gamma_\mu q_6 b | \Lambda_b \rangle$) transition is parametrized in terms of three form factors $f_0^T$, $f_0^T$, $f_0^T$ ($f_0^T$, $f_0^T$, $f_0^T$). Therefore, fourteen $q^2$ dependent form factors contribute to this decay.

With the parametrizations of $\Lambda \to \Lambda^*$ transitions at our disposal, we now calculate the hadronic amplitudes. In literature, instead of the helicity amplitudes, the so called transversity amplitudes are often used. The transversity amplitudes are linear combinations of the helicity amplitudes, see Appendix E. For the (axial-)vectors currents they read

$$B_{\perp 1}^{L(R)} = \sqrt{2} N \left( f_0^T \sqrt{s_{\Lambda b}} C_{VA}^{L(R)} + \frac{2m_b}{q^2} f_0^T \sqrt{s_{\perp}} (C_7 + C_7') \right),$$ 
(2.11)

$$B_{\parallel 1}^{L(R)} = \sqrt{2} N \left( f_0^A \sqrt{s_{\Lambda b}} C_{VA}^{L(R)} + \frac{2m_b}{q^2} f_0^T \sqrt{s_{\perp}} (C_7 - C_7') \right),$$ 
(2.12)

$$B_{\perp 0}^{L(R)} = -\sqrt{2} N \left( f_0^T \frac{(m_{\Lambda_b} + m_{\Lambda^*}) s - \sqrt{s_{\perp}}}{\sqrt{6m_{\Lambda^*}}} C_{VA}^{L(R)} + \frac{2m_b}{q^2} f_0^T \sqrt{s_{\perp}} \frac{s - \sqrt{s_{\perp}}}{\sqrt{6m_{\Lambda^*}}}, (C_7 + C_7') \right),$$ 
(2.13)

$$B_{\parallel 0}^{L(R)} = \sqrt{2} N \left( f_0^A \frac{(m_{\Lambda_b} - m_{\Lambda^*}) s + \sqrt{s_{\perp}}}{\sqrt{6m_{\Lambda^*}}} C_{VA}^{L(R)} + \frac{2m_b}{q^2} f_0^T \sqrt{s_{\perp}} \frac{s + \sqrt{s_{\perp}}}{\sqrt{6m_{\Lambda^*}}}, (C_7 - C_7') \right),$$ 
(2.14)

$$A_{\perp 1}^{L(R)} = -\sqrt{2} N \left( f_1^T \frac{s - \sqrt{s_{\perp}}}{\sqrt{3m_{\Lambda^*}}} C_{VA}^{L(R)} + \frac{2m_b}{q^2} f_1^T \frac{s + \sqrt{s_{\perp}}}{\sqrt{3m_{\Lambda^*}}}, (C_7 + C_7') \right),$$ 
(2.15)
\[ A_{\parallel 1}^{L(R)} = -\sqrt{2N} f_t \left( \lambda^{s+\sqrt{s}} C_{VA+} + \frac{2m_b}{q^2} f_t f_1^5 (m_{\Lambda_b} - m_{\Lambda^*}) \frac{s+\sqrt{s}}{\sqrt{3m_{\Lambda^*}}} (C_7 - C'_7) \right), \]  
\[ A_{\perp t}^{L(R)} = \sqrt{2N} f_t \left( \frac{m_{\Lambda_b} - m_{\Lambda^*}}{m_b - m_s} \frac{s+\sqrt{s}}{\sqrt{6m_{\Lambda^*}}} C_{SP+} \right), \]  
\[ A_{\parallel t}^{L(R)} = -\sqrt{2N} f_t \left( \frac{m_{\Lambda_b} + m_{\Lambda^*}}{m_b + m_s} \frac{s-\sqrt{s}}{\sqrt{6m_{\Lambda^*}}} C_{SP-} \right), \]

where
\[ s_\pm = (m_{\Lambda_b} \pm m_{\Lambda^*})^2 - q^2, \]  
and the Wilson coefficients \( c_{VA\pm}^{L(R)} \) are defined as follows
\[ c_{VA+}^{L(R)} = (C_9 \mp 10) + (C_{9'} \mp 10'), \]  
\[ c_{VA-}^{L(R)} = (C_9 \mp 10) - (C_{9'} \mp 10'). \]  

For the (pseudo-)scalar currents the amplitudes read
\[ A_{S\perp}^{L(R)} = \sqrt{2N} f_t \left( \frac{m_{\Lambda_b} - m_{\Lambda^*}}{m_b - m_s} \frac{s+\sqrt{s}}{\sqrt{6m_{\Lambda^*}}} C_{SP+} \right), \]  
\[ A_{S\parallel}^{L(R)} = -\sqrt{2N} f_t \left( \frac{m_{\Lambda_b} + m_{\Lambda^*}}{m_b + m_s} \frac{s-\sqrt{s}}{\sqrt{6m_{\Lambda^*}}} C_{SP-} \right), \]  

where the scalar Wilson coefficients are defined as
\[ c_{SP+}^{L(R)} = (C_S + C_{S'}) \mp (C_P + C_{P'}), \]  
\[ c_{SP-}^{L(R)} = (C_S - C_{S'}) \mp (C_P - C_{P'}). \]  

The overall normalization factor that has been customarily absorbed in the transversity amplitude is given by
\[ N = G_F V_{tb} V_{ts}^* \alpha_{\perp} \sqrt{\frac{q^2 \sqrt{\lambda (m_{\Lambda_b}^2, m_{\Lambda^*}^2, q^2)}}{3 \cdot 211 m_{\Lambda_b}^3 \pi^5} \beta_{\ell} B_{\Lambda^*}}, \]  
\[ \beta_{\ell} = \sqrt{1 - \frac{4m_{\ell}^2}{q^2}}, \]

where \( B_{\Lambda^*} = B_{\Lambda^*}(\Lambda^* \to N\bar{K}) \) is the branching ratio and \( \tau_{\Lambda_b} \) is the \( \Lambda_b \) lifetime.

### 2.4 The \( \Lambda^* \to N\bar{K} \) decay

The effective Lagrangian describing the strong decay \( \Lambda^* \to N\bar{K} \) is assumed to be \(^1\) [34]
\[ \mathcal{L}_1 = g m_{\Lambda^*} \bar{\psi}_\mu (g^{\mu\nu} + \alpha \gamma^\mu \gamma^\nu) \gamma_5 \Psi \partial_\nu \phi + h.c., \]  

where \( g \) is a coupling constant, \( \Psi \) is a spin-1/2 field describing the \( \Lambda_b \), and \( \phi \) is scalar field corresponding to the \( \bar{K} \) meson. The \( \Lambda^* \) is a spin 3/2 particle and is described by a Rarita-Schwinger field \( \psi_\mu \) [36]. In addition to the vector index, there is an implicit spinor

\(^1\) A different choice for the Lagrangian is given in [35] which leads to same result.
index in the Rarita-Schwinger field. The parameter $a$ is only relevant for loop calculations. The Hamiltonian (2.25) yields the following matrix element for $\Lambda^* \to N\bar{K}$ transition

$$\mathcal{M}_{\Lambda^*}(s_{\Lambda^*}, s_N) = g m_{\Lambda^*} k_2^\mu u^{s_N}_\gamma \gamma_5 U^{a\Lambda^*}_\mu,$$  \hspace{1cm} (2.26)

where $U^{a\Lambda^*}_\mu$ is the Rarita-Schwinger spinor describing the $\Lambda^*$ and $u^{s_N}$ is the Dirac spinor for the $N$. In the rest frame of the $\Lambda^*$ the solutions of Rarita-Schwinger and Dirac spinors are given in Appendix F. As can be understood from (2.5), the matrix elements $\mathcal{M}_{\Lambda^*}$ contribute to the $\Lambda_b \to \Lambda^*(\to N\bar{K})\ell^+\ell^-$ through the following interference term

$$\Gamma_2(s^a_{\Lambda^*}, s^b_{\Lambda^*}) = \frac{\sqrt{T_T - r_T}}{16 m_{\Lambda^*} \pi^3} \sum s_N \mathcal{M}_{\Lambda^*}(s^a_{\Lambda^*}, s_N)[\mathcal{M}_{\Lambda^*}(s^b_{\Lambda^*}, s_N)]^*, \hspace{1cm} (2.27)$$

where

$$r_T = (m_{\Lambda^*} \pm m_N)^2 - m_K^2. \hspace{1cm} (2.28)$$

Using the solutions of Rarita-Schwinger and Dirac fields given in Appendix. F we obtain \(^2\)

$$\Gamma_2(s^a_{\Lambda^*}, s^b_{\Lambda^*}) = \frac{\Gamma(\Lambda^* \to N\bar{K})}{4} \times \begin{pmatrix} 6\sin^2(\theta_{\Lambda^*}) & -2\sqrt{3}e^{-i\phi}\sin(2\theta_{\Lambda^*}) & -2\sqrt{3}e^{-2i\phi}\sin^2(\theta_{\Lambda^*}) & 0 \\
-2\sqrt{3}e^{i\phi}\sin(2\theta_{\Lambda^*}) & 3\cos(2\theta_{\Lambda^*}) + 5 & 0 & -2\sqrt{3}e^{-2i\phi}\sin^2(\theta_{\Lambda^*}) \\
-2\sqrt{3}e^{2i\phi}\sin^2(\theta_{\Lambda^*}) & 0 & 3\cos(2\theta_{\Lambda^*}) + 5 & +2\sqrt{3}e^{-i\phi}\sin(2\theta_{\Lambda^*}) \\
0 & -2\sqrt{3}e^{2i\phi}\sin^2(\theta_{\Lambda^*}) & +2\sqrt{3}e^{i\phi}\sin(2\theta_{\Lambda^*}) & 6\sin^2(\theta_{\Lambda^*}) \end{pmatrix}, \hspace{1cm} (2.29)$$

where the $\Lambda^* \to N\bar{K}$ decay width is defined as

$$\Gamma(\Lambda^* \to N\bar{K}) = \frac{1}{4} \sum s^a_{\Lambda^*} \Gamma_2(s^a_{\Lambda^*}, s^b_{\Lambda^*}). \hspace{1cm} (2.30)$$

### 3 Angular distributions

The results of the previous sections yield the following four fold angular distributions for $\Lambda_b \to \Lambda^*(\to N\bar{K})\ell^+\ell^-$ decay

$$\frac{d^4B}{dq^2dcos \theta_\ell dcos \theta_{\Lambda^*}d\phi} = \frac{3}{8\pi} \left[ (K_{1c}\cos \theta_\ell + K_{1cc}\cos^2 \theta_\ell + K_{1ss}\sin^2 \theta_\ell)\cos^2 \theta_{\Lambda^*} \\
+ (K_{2c}\cos \theta_\ell + K_{2cc}\cos^2 \theta_\ell + K_{2ss}\sin^2 \theta_\ell)\sin^2 \theta_{\Lambda^*} \\
+ (K_{3ss}\sin^2 \theta_\ell)\sin^2 \theta_{\Lambda^*}\cos \phi + (K_{4ss}\sin^2 \theta_\ell)\sin^2 \theta_{\Lambda^*}\sin \phi \cos \phi \\
+ (K_{5ss}\sin \theta_\ell + K_{5sc}\sin \theta_\ell\cos \theta_\ell)\sin \theta_{\Lambda^*}\cos \theta_{\Lambda^*}\cos \phi \\
+ (K_{6ss}\sin \theta_\ell + K_{6sc}\sin \theta_\ell\cos \theta_\ell)\sin \theta_{\Lambda^*}\cos \theta_{\Lambda^*}\sin \phi \cos \phi \right]. \hspace{1cm} (3.1)$$

\(^2\)We differ with Ref.[17] by sign of some of the entries of the $4 \times 4$ matrix.
The $K_{\{\ldots\}}$, where $\{\ldots\} = 1c, \ldots 6sc$, are called the angular coefficients that can be written in terms of the transversity amplitudes. As the masses of the final states has been kept, we show the mass corrections of the order $O(m_\ell/\sqrt{q^2})$ and $O(m_\ell^2/q^2)$ and write the angular coefficients as

$$K_{\{\ldots\}} = \mathcal{K}_{\{\ldots\}} + \frac{m_\ell}{\sqrt{q^2}} K'_{\{\ldots\}} + \frac{m_\ell^2}{q^2} K''_{\{\ldots\}}.$$  \hfill (3.2)

The detailed expressions of $K_{\{\ldots\}}, K'_{\{\ldots\}}$ and $K''_{\{\ldots\}}$ in terms of the transversity amplitudes are given in the Appendix G. The terms $K'$ and $K''$ are important if the final state leptons are significantly heavy, for example $\tau$ leptons.

4 Observables

From the four-fold differential distribution (3.1) we construct observables that are suitable for experimental analysis. By weighted angular integrals of the differential distribution we obtain different observables as a liner combinations of the angular coefficients and as a function of $q^2$

$$O[\omega] = \int \frac{d^4B}{dq^2 d\cos \theta_\ell d\cos \theta_\Lambda^* d\phi} \omega(q^2, \theta_\ell, \theta_\Lambda^*, \phi) d\cos \theta_\ell d\cos \theta_\Lambda^* d\phi.$$  \hfill (4.1)

We limit ourselves to the following ‘simple’ observables

- $\omega = 1$ yields the simplest observable, differential branching ratio as a function of $q^2$

$$\frac{d\mathcal{B}}{dq^2} = \frac{2}{3} \left[ K_{1cc} + 2K_{1ss} + 2K_{2ss} + 4K_{2ss} + 2K_{3ss} \right].$$  \hfill (4.2)

- The fraction of longitudinal polarization of the dilepton system

$$F_L = 1 - \frac{2(K_{1cc} + 2K_{2cc})}{K_{1cc} + 2(K_{1ss} + K_{2cc} + 2K_{2ss} + K_{3ss})},$$  \hfill (4.3)

is obtained for the choice

$$\omega = \frac{2 - 5 \cos^2 \theta_\ell}{d\mathcal{B}/dq^2}.$$  

- The choice

$$\omega = \frac{\text{sgn} \{\cos \theta_\ell\}}{d\mathcal{B}/dq^2},$$  \hfill (4.4)

yields the well known lepton forward-backward asymmetry

$$A_\ell^{FB} = \frac{3(K_{1c} + 2K_{2c})}{2(K_{1cc} + 2(K_{1ss} + K_{2cc} + 2K_{2ss} + K_{3ss}))}.$$  \hfill (4.5)

Since the $\Lambda^*$ decays through strong interaction, the hadronic side asymmetry $A_{\Lambda^*}^{FB}$ and the mixed asymmetry $A_{\Lambda^*}^{FB}$ [16] vanish.
5 \( \Lambda_b \rightarrow \Lambda^* \) in HQET

The description of \( \Lambda_b \rightarrow \Lambda^* \) transition involves fourteen \( q^2 \)-dependent form factors. In the limit of the low hadronic recoil, the number of independent form factors reduce as a consequence of the HQET spin symmetry \([24–26, 31, 37]\). In this section we derive the improved Isgur-Wise relations between the form factors to the leading order in \( 1/m_b \) and including \( \mathcal{O}(\alpha_s) \) corrections.

5.1 Improved Isgur-Wise relations

The starting point is to use the heavy baryon velocity \( v^\mu = p^\mu/m_{\Lambda_b} \) to project the \( b \)-quark filed on to its large spinor components \( h_v = \not{v} h_v \) in terms of leading Isgur-Wise form factors \( \xi_{1,2} \)

\[
\langle \Lambda^*(k, s_{\Lambda^*}) | s\Gamma b \Lambda_b(p = \not{v} m_{\Lambda_b}, s_{\Lambda_b}) \rangle \rightarrow \langle \Lambda^*(k, s_{\Lambda^*}) | s\Gamma h_v | \Lambda_b(v, s_{\Lambda_b}) \rangle \\
\simeq \bar{U}^\alpha_{\Lambda^*}(k, s_{\Lambda^*}) v_\alpha (\xi_1 + \frac{\xi_2}{v}) \Gamma u_{\Lambda_b}(v, s_{\Lambda_b}), \quad (5.1)
\]

where \( U^\alpha_{\Lambda^*} \) is the Rarita-Schwinger spinor describing the \( \Lambda^* \), \( u_{\Lambda_b} \) is the usual Dirac spinor describing the \( \Lambda_b \), and \( \Gamma \) is a Dirac matrix. The Isgur-Wise form factors are defined as

\[
\xi_{1,2} \equiv \xi_{1,2}(v.k).
\]

In the limit of heavy quark \( m_b \rightarrow \infty \), i.e., neglecting the contributions of the order \( 1/m_b \) in the definitions of the form factors \( (D.1), (D.6) \) and comparing with \( (5.1) \) we get

\[
\begin{align*}
\frac{f_V}{f_A} &= \frac{f_0^V}{f_0^A} = \frac{f_T^V}{f_T^A} = f_0^T = \frac{\xi_1 - \xi_2}{m_{\Lambda_b}}, \\
\frac{f_A}{f_V} &= \frac{f_0^A}{f_0^V} = \frac{f_T^A}{f_T^V} = f_0^T = \frac{\xi_1 + \xi_2}{m_{\Lambda_b}}, \\
\frac{f_A}{f_T} &= \frac{f_0^A}{f_0^T} = \frac{f_5^T}{f_5^A} = f_5^T = 0.
\end{align*}
\]

To account for the sub-leading corrections of the order \( \mathcal{O}(\alpha_s) \), following \[31\] we use the matching relations of the QCD currents onto the HQET. For the (axial-)vector currents the relation is

\[
\bar{s}\gamma^\mu b = C_0^{(v)} \bar{s}\gamma^\mu h_v + C_1^{(v)} v^\mu \bar{s}h_v + \frac{1}{2m_b} \bar{s}\gamma^\mu i\not{\nabla}_\perp h_v + \cdots,
\]

\[
\bar{s}\gamma^\mu \gamma_5 b = C_0^{(v)} \bar{s}\gamma^\mu \gamma_5 h_v - C_1^{(v)} v^\mu \bar{s}\gamma_5 h_v - \frac{1}{2m_b} \bar{s}\gamma^\mu i\not{\nabla}_\perp \gamma_5 h_v + \cdots,
\]

and for the (pseudo-)tensor currents the relation is

\[
\bar{s}i\sigma^{\mu\nu} q_\nu(\gamma_5)b = C_0^{(t)} \bar{s}i\sigma^{\mu\nu} q_\nu(\gamma_5) h_v \pm \frac{1}{2m_b} \bar{s}\sigma^{\mu\nu} q_\nu i\not{\nabla}_\perp (\gamma_5) h_v + \cdots.
\]
The renormalization scale $\mu$-dependent matching coefficients $C_{0}^{(v,t)}, C_{1}^{(v)}$ at next-to-leading order in $\alpha_s$ are [31]

\begin{align}
C_{0}^{(v)} &= 1 - \frac{\alpha_s C_F}{4\pi} \left( 3 \ln \left( \frac{\mu}{m_b} \right) + 4 \right) + O(\alpha_s^2), \\
C_{1}^{(v)} &= \frac{\alpha_s C_F}{2\pi} + O(\alpha_s^2), \\
C_{0}^{(t)} &= 1 - \frac{\alpha_s C_F}{4\pi} \left( 5 \ln \left( \frac{\mu}{m_b} \right) + 4 \right) + O(\alpha_s^2) .
\end{align}

The matrix elements of these currents can be parametrized in terms of the leading Isgur-Wise form factors $\xi_{1,2}$ as

\begin{align}
\langle \Lambda^*(k, s_{\Lambda'}) | \bar{s} \gamma^\mu (\gamma_5) b | \Lambda_b(p = m_{\Lambda_b} v, s_{\Lambda_b}) \rangle \\
&\simeq C_{0}^{(v)} \sum_{n=1,2} \xi_n U_{\Lambda^*}^\alpha (k, s_{\Lambda'}) v_n \Gamma_n \gamma^\mu (\gamma_5) u_{\Lambda_b}(v, s_{\Lambda_b}) \\
&\pm C_{1}^{(v)} \sum_{n=1,2} \xi_n v_\mu U_{\Lambda^*}^\alpha (k, s_{\Lambda'}) v_n \Gamma_n (\gamma_5) u_{\Lambda_b}(v, s_{\Lambda_b}), \\
\langle \Lambda^*(k, s_{\Lambda'}) | \bar{s} \gamma^\mu \gamma^\nu q_\nu (\gamma_5) b | \Lambda_b(p = m_{\Lambda_b} v, s_{\Lambda_b}) \rangle \\
&\simeq C_{0}^{(t)} \sum_{n=1,2} \xi_n U_{\Lambda^*}^\alpha (k, s_{\Lambda'}) v_n \Gamma_n i\sigma^{\mu\nu} q_\nu (\gamma_5) u_{\Lambda_b}(v, s_{\Lambda_b}),
\end{align}

where the two independent Dirac structures are

\begin{align}
\Gamma_1 = 1, \quad \Gamma_2 = i. \tag{5.14}
\end{align}

Comparing the parametrizations (5.12) and (5.13) with (D.1)-(D.6) we get the following expressions for the physical form factors at leading order in $1/m_b$

\begin{align}
f_{V}^{A} &= C_{0}^{(v)} \frac{\xi_1 + \xi_2}{m_{\Lambda_b}}, \\
\bar{f}_{0}^{V} &= C_{0}^{(v)} \left( \frac{C_{1}^{(v)} s_{\pm}}{2m_{\Lambda_b}(m_{\Lambda_b} \pm m_{\Lambda^*})} \right) \frac{\xi_1}{m_{\Lambda_b}} \mp \left( C_{0}^{(v)} - \frac{2C_{0}^{(v)} + C_{1}^{(v)} s_{\pm}}{2m_{\Lambda_b}(m_{\Lambda_b} \pm m_{\Lambda^*})} \right) \frac{\xi_2}{m_{\Lambda_b}}, \\
f_{T}^{(5)} &= \frac{C_{0}^{(v)}}{m_{\Lambda_b}} \left( \frac{\xi_1 + \xi_2}{m_{\Lambda_b}} \mp \frac{s_{\pm}}{m_{\Lambda_b}(m_{\Lambda_b} \pm m_{\Lambda^*})} \right), \\
\bar{f}_{0}^{T} &= C_{0}^{(v)} \frac{\xi_1 \mp \xi_2}{m_{\Lambda_b}}, \\
f_{V}^{(q^2)} &= \frac{1}{m_{\Lambda_b}} \xi_1 \left( C_{0}^{(v)} + C_{1}^{(v)} \left( 1 - \frac{s_{-}}{2m_{\Lambda_b}(m_{\Lambda_b} - m_{\Lambda^*})} \right) \right) \\
&\quad + \frac{1}{m_{\Lambda_b}} \xi_2 \left( C_{0}^{(v)} \left( 1 - \frac{s_{-}}{m_{\Lambda_b}(m_{\Lambda_b} - m_{\Lambda^*})} \right) + C_{1}^{(v)} \left( 1 - \frac{s_{-}}{2m_{\Lambda_b}(m_{\Lambda_b} - m_{\Lambda^*})} \right) \right), \\
\bar{f}_{V}^{(q^2)} &= \frac{1}{m_{\Lambda_b}} \xi_1 \left( C_{0}^{(v)} + C_{1}^{(v)} \left( 1 - \frac{s_{+}}{2m_{\Lambda_b}(m_{\Lambda_b} + m_{\Lambda^*})} \right) \right) \\
&\quad - \frac{1}{m_{\Lambda_b}} \xi_2 \left( C_{0}^{(v)} \left( 1 - \frac{s_{+}}{m_{\Lambda_b}(m_{\Lambda_b} + m_{\Lambda^*})} \right) + C_{1}^{(v)} \left( 1 - \frac{s_{+}}{2m_{\Lambda_b}(m_{\Lambda_b} + m_{\Lambda^*})} \right) \right).
\end{align}
The form factors $f_{g}^{V,A}$ remain zero. These expressions will be used to correlate the form factors and reduce the number of independent form factors in the transversity amplitudes.

### 5.2 Low-recoil factorization

The improved Isgur-Wise relations (5.15)-(5.20) lead to simplifications of the description of the decay in the low recoil region. In what follows, we consider the (axial-)vector form factors $f_{0}^{V,A}$, $f_{\perp}^{V,A}$, $f_{t}^{V,A}$ as independent and use the improved Isgur-Wise relations (5.15)-(5.20) to relate the tensor and pseudo-tensor form factors. With this consideration, at leading order in $1/m_\text{b}$ and including $\mathcal{O}(\alpha_s)$ corrections, the (axial-)vector type transversity amplitudes $B_{\perp 1}^{(L,R)}$, $B_{\parallel 1}^{(L,R)}$ vanish, and $A_{\perp 1,0}^{(L,R)}$ and $A_{\parallel 1,0}^{(L,R)}$ depend only on a single form factors

$$B_{\perp 1}^{(L,R)} = 0, \quad B_{\parallel 1}^{(L,R)} = 0,$$

$$A_{\perp 0}^{(L,R)} = -\frac{2N^\pm m_{\Lambda_b} + m_{\Lambda^*}}{\sqrt{q^2}} \frac{s-\sqrt{s_\perp}}{\sqrt{6m_{\Lambda^*}}} C_+^{L(R)} f_0^V,$$

$$A_{\parallel 0}^{(L,R)} = \frac{2N^\pm m_{\Lambda_b} - m_{\Lambda^*}}{\sqrt{q^2}} \frac{s+\sqrt{s_\perp}}{\sqrt{6m_{\Lambda^*}}} C_-^{L(R)} f_0^A,$$

$$A_{\perp 1}^{(L,R)} = \frac{N_s^\pm \sqrt{s_\perp}}{\sqrt{3m_{\Lambda^*}}} C_{\pm}^{L(R)} f_\perp^V,$$

$$A_{\parallel 1}^{(L,R)} = \frac{N_s^\pm \sqrt{s_\perp}}{\sqrt{3m_{\Lambda^*}}} C_{\pm}^{L(R)} f_\parallel^A,$$

where the combinations of Wilson coefficients are

$$C_+^{L(R)} = (C_9 + C_9') + \frac{2\kappa m_\text{b} m_{\Lambda_b}}{q^2} (C_7 + C_7') \mp (C_{10} + C_{10'}),$$

$$C_-^{L(R)} = (C_9 - C_9') + \frac{2\kappa m_\text{b} m_{\Lambda_b}}{q^2} (C_7 - C_7') \mp (C_{10} - C_{10'}).$$

The parameter

$$\kappa \equiv \kappa(\mu) = \frac{C_{0}^{(t)}}{C_{0}^{(v)}} = 1 - \frac{\alpha_s C_F}{2\pi} \ln \left( \frac{\mu}{m_\text{b}} \right),$$

absorbs the perturbative QCD corrections to the form factor relations in such a way that with the product of the Wilson coefficients and the $b$-quark mass, the transversity amplitudes are free of renormalization-scale at a given order in the perturbation theory. In this derivation, we have ignored the sub-leading terms of the order $m_{\Lambda^*}/m_{\Lambda_b}$ and $\Lambda_{\text{QCD}}/m_{\Lambda_b}$, and used a naively anti-commuting $\gamma_5$ matrix. The consequences of the simplified expressions of the transversity amplitudes translate to factorization of short- and long-distance physics in the decay observables. The factorization in scale makes the electroweak physics transparent in the observables. In the subsequent subsections the low-recoil factorization is discussed. For these discussions we will neglect the mass of the leptons which is a valid approximation in the low-recoil region.
5.2.1 In SM basis

In the SM+SM’ basis there are four distinct Wilson coefficients $C^{L(R)}_{\pm}$ Eq. (5.26). In the observables the Wilson coefficients appear as following short-distance coefficients

$$\rho_1^+ = \frac{1}{2} \left( |C^R_\perp|^2 + |C^L_\perp|^2 \right), \quad \rho_2^+ = \frac{1}{4} \left( C^R_\perp C^R_\perp^\ast + C^L_\perp C^L_\perp^\ast \right).$$  \hspace{1cm} (5.28)

The coefficients $\rho_1^+$ and $\rho_2^+$ are present in $B \to K^* \ell^+ \ell^-$ [38], non-resonant $B \to K \pi \ell^+ \ell^-$ [39], and $\rho_2^-$ appear in $\Lambda_b \to \Lambda(\to N \pi) \ell^+ \ell^-$ [11] decay. In the SM $C_{9(10),10'} = 0$ and $C_{S(1),F(2)} = 0$, and only two Wilson coefficients are non-vanishing

$$C^{L(R)}_{\perp} \equiv C^{L(R)}_{\perp} = C^{L(R)}_9 + \frac{2\kappa M_b m_{\Lambda_b}}{q^2} C_7 + C_{10}. \hspace{1cm} (5.29)$$

Therefore, in the SM observables only two short-distance coefficients $\rho_{1,2}$ are relevant [39]

$$\rho_1 \equiv \rho_1^+ = 2\text{Re} \rho_2^+ , \quad \rho_2 \equiv \text{Re} \rho_2^+ , \quad \text{Im} \rho_2^+ = 0. \hspace{1cm} (5.30)$$

The low recoil factorization of the angular coefficients read

$$K_{1c} = \frac{8N^2}{3m_{\Lambda^+}^2} (s_+ s_-)^{3/2} f^A_\perp f^A_\perp \rho_2,$$  \hspace{1cm} (5.31)

$$K_{1cc} = \frac{2N^2}{3m_{\Lambda^+}^2} s_+ s_- \left( s_- |f^A_\perp|^2 + s_+ |f^A_\perp|^2 \right) \rho_1,$$  \hspace{1cm} (5.32)

$$K_{1ss} = \frac{N^2}{3m_{\Lambda^+}^2} s_+ s_- \left( \frac{(m_{\Lambda_b} + m_{\Lambda^+})^2}{q^2} s_- |f^A_0|^2 + \frac{(m_{\Lambda_b} - m_{\Lambda^+})^2}{q^2} s_+ |f^A_0|^2 \right.$$  
$$+ s_- |f^A_\perp|^2 + s_+ |f^A_\perp|^2 \right) \rho_1,$$  \hspace{1cm} (5.33)

$$K_{2c} = \frac{2N^2}{3m_{\Lambda^+}^2} (s_+ s_-)^{3/2} f^A_\perp f^A_\perp \rho_2,$$  \hspace{1cm} (5.34)

$$K_{2cc} = \frac{N^2}{6m_{\Lambda^+}^2} s_+ s_- \left( s_- |f^A_\perp|^2 + s_+ |f^A_\perp|^2 \right) \rho_1,$$  \hspace{1cm} (5.35)

$$K_{2ss} = \frac{N^2}{12m_{\Lambda^+}^2} s_+ s_- \left( \frac{(m_{\Lambda_b} + m_{\Lambda^+})^2}{q^2} s_- |f^A_0|^2 + \frac{(m_{\Lambda_b} - m_{\Lambda^+})^2}{q^2} s_+ |f^A_0|^2 \right.$$  
$$+ s_- |f^A_\perp|^2 + s_+ |f^A_\perp|^2 \right) \rho_1.$$  \hspace{1cm} (5.36)

The simple observables $dB/dq^2$, $A^L_{FB}$ and $F_L$ read

$$\frac{dB}{dq^2} = \frac{N^2 s_+ s_-}{3m_{\Lambda^+}^2} \left[ s_- \left( 2 |f^A_\perp|^2 + \frac{(m_{\Lambda_b} + m_{\Lambda^+})^2}{q^2} |f^A_0|^2 \right) \right.$$  
$$+ s_+ \left( 2 |f^A_\perp|^2 + \frac{(m_{\Lambda_b} + m_{\Lambda^+})^2}{q^2} |f^A_0|^2 \right) \right] \rho_1,$$  \hspace{1cm} (5.37)

$$\frac{dB}{dq^2} F_L = \frac{dB}{dq^2} - \frac{2N^2}{3m_{\Lambda^+}^2} s_+ s_- \left( s_- |f^A_\perp|^2 + s_+ |f^A_\perp|^2 \right) \rho_1,$$  \hspace{1cm} (5.38)

$$\frac{dB}{dq^2} A^L_{FB} = \frac{2N^2}{m_{\Lambda^+}^2} (s_+ s_-)^{3/2} f^A_\perp f^A_\perp \rho_2.$$  \hspace{1cm} (5.39)
The goal of the future experiments would be to test the electroweak physics by extracting the short-distance coefficients and to simultaneously extract the form factors. In this regard, it is useful to discuss some observables where for example the form factors cancels and the short-distance coefficients can be easily extracted. Following three ratios are constants

\[ \frac{K_{1c}}{K_{2c}} = 4, \quad \frac{K_{1cc}}{K_{2cc}} = 4, \quad \frac{K_{1ss}}{K_{2ss}} = 4, \quad (5.40) \]

and the following ratios depends on only one ratio of form factor and \( \rho_1/\rho_2 \)

\[ \frac{K_{1c}}{K_{1cc}} = \frac{K_{2c}}{K_{2cc}} = \left( \frac{4\sqrt{s_+s_-}}{s_+ f^2_{t^A} + s_+ f^2_{t^A}} \right) \rho_2 \rho_1, \quad (5.41) \]

The ratio is suitable to test the short distance physics. The following ratios are independent of any short-distance physics

\[ \frac{K_{1ss}}{K_{1cc}} = \frac{K_{2ss}}{K_{2cc}} = \frac{1}{2} \left[ 1 + \frac{(m_{A_b} + m_{A_c})^2 s_- |f^V_0|^2 + (m_{A_b} - m_{A_c})^2 s_+ |f^A_0|^2}{q^2 (s_+ f^2_{t^A} + s_+ f^2_{t^A})} \right]. \quad (5.42) \]

If the new physics is absent, the ratio \( K_{1ss}/K_{1cc} \) is suitable to test the form factors.

5.2.2 In SM+SM'+SP basis

When the SM operator basis in extended by the SM+SM'+SP operator basis we obtain

\[ K_{1c} = \frac{8N^2}{3m_{A^*}} (s_+ s_-)^{3/2} f^V_{t^A} f^A_{t^A} \rho_2^+, \quad (5.43) \]

\[ K_{1cc} = \frac{2N^2}{3m_{A^*}} s_+ s_- \left( s_- \rho_1^+ |f^V_{t^A}|^2 + s_+ \rho_1^- |f^A_{t^A}|^2 \right) + \frac{1}{2} \left( \frac{(m_{A_b} - m_{A_c})^2}{(m_b - m_a)^2} s_+ \rho_2^+ |f^V_{t^A}|^2 + \frac{(m_{A_b} + m_{A_c})^2}{(m_b + m_a)^2} s_- \rho_2^- |f^A_{t^A}|^2 \right), \quad (5.44) \]

\[ K_{1ss} = \frac{N^2}{3m_{A^*}} s_+ s_- \left( \frac{(m_{A_b} + m_{A_c})^2}{(m_b + m_a)^2} s_- \rho_1^+ |f^V_{t^A}|^2 + \frac{(m_{A_b} - m_{A_c})^2}{(m_b - m_a)^2} s_+ \rho_1^- |f^A_{t^A}|^2 \right) \]

\[ + \frac{(m_{A_b} + m_{A_c})^2}{(m_b + m_a)^2} s_+ \rho_2^+ |f^V_{t^A}|^2 - \frac{(m_{A_b} + m_{A_c})^2}{(m_b - m_a)^2} s_+ \rho_2^- |f^A_{t^A}|^2 \], \quad (5.45) \]

\[ K_{2c} = \frac{2N^2}{3m_{A^*}} (s_+ s_-)^{3/2} f^V_{t^A} f^A_{t^A} \rho_2^+, \quad (5.46) \]

\[ K_{2cc} = \frac{N^2}{6m_{A^*}} s_+ s_- \left( s_- \rho_1^+ |f^V_{t^A}|^2 + s_+ \rho_1^- |f^A_{t^A}|^2 \right) + \frac{4}{3} \left( \frac{(m_{A_b} - m_{A_c})^2}{(m_b - m_a)^2} s_+ \rho_2^+ |f^V_{t^A}|^2 + \frac{(m_{A_b} + m_{A_c})^2}{(m_b + m_a)^2} s_- \rho_2^- |f^A_{t^A}|^2 \right), \quad (5.47) \]

\[ K_{2ss} = \frac{N^2}{12m_{A^*}} s_+ s_- \left( \frac{(m_{A_b} + m_{A_c})^2}{q^2} s_- \rho_1^+ |f^V_{t^A}|^2 + \frac{(m_{A_b} - m_{A_c})^2}{q^2} s_+ \rho_1^- |f^A_{t^A}|^2 \right). \]
The rest of the coefficients vanish in this operator basis. Here we have defined
\[ \rho_{S}^{\pm} = \frac{1}{2} \left( |C_{SP+}^{R}|^2 + |C_{SP+}^{L}|^2 \right). \] (5.49)
The list of input parameters are listed in the table 1.

We make the following observations

- The ratios \( K_{1c}/K_{2c} \) and \( K_{1ss}/K_{2ss} \) remain independent of both short- and long-distance physics in the most general SM+SM’+SP basis.
- The ratio \( K_{1cc}/K_{2cc} \) remains independent of short- and long-distance physics in the SM+SM’ basis. If the SP basis is present then it depends on both form factor and electroweak physics.
- The ratios \( K_{1ss}/K_{1cc} \) and \( K_{1c}/K_{1cc} \) both are sensitive to NP.

6 Numerical Analysis

In this section we perform a numerical analysis of the \( \Lambda_{b} \rightarrow \Lambda^{*} (\rightarrow N\bar{K}) \mu^{+}\mu^{-} \) observables and study their sensitivity to the SM and NP. The requisite inputs are the form factors. The presently available lattice QCD determination of the form factors are only preliminary [40] and we use the estimates from the non-relativistic quark model [41]. Since the Ref. [41] does not give an estimate of the uncertainties of the form factors, we assume uncorrelated 30% uncertainties on them for illustrative purpose. Since the form factor estimates are crude, we refrain from doing an intricate error analysis involving other parametric inputs. The list of input parameters are listed in the table 1.
As mentioned in the Sec. 2 the short-distance inputs $C_{T,9}$ include the factorizable charm-loop contributions coming from the hadronic operators $O_{1-6}$ and $O_8$ [31]. The SM values of all the Wilson coefficients are taken from [42], evaluated at the $b$-quark mass scale $\mu = m_b = 4.8$ GeV. The $b$- and the $c$-quark loop functions are taken from [22, 43–45].

In order to ascertain the effects of the NP coefficients, we first discuss the model-independent constraints on the Wilson coefficients. For the simplicity of the discussions we consider two scenarios-- (i) only the chirality flipped NP (SM+SM$'$ basis) is present, and (ii) only scalar NP is present (SM+SP basis). For scenario (i) model-independent constraints on $\delta C_{9,10}^{NP}$, $\delta C_{9,10}'^{NP}$ are available from the global fits to $b \to s \mu^+\mu^-$ data. For demonstration we restrict ourselves to the following benchmark solutions [46, 47]

$$\delta C_9^{NP} = -1.11, \quad \delta C_9'^{NP} = -1.01, \quad \delta C_{10}^{NP} = -1.16, \quad \delta C_{10}'^{NP} = 0.38. \quad (6.1)$$

In scenario (ii) model-independent constraints on $C_{S(P)}^{(i),P}$ come from the experimental data on $B_s \to \mu^+\mu^-$ and the inclusive $B \to X_s \mu^+\mu^-$ decays [12]

$$C_{S(P)}^{(i),P} = [-4.0, +4.0]. \quad (6.4)$$

To the best of our knowledge there is no model-independent global fit that considers a scenario where both chirality flipped and scalar operators are simultaneously present. This scenario is therefore not discussed in this paper.

In Fig. 1 we have shown the SM estimates of the observables $dB/dq^2$, $A_{FB}^\ell$ and $F_L$ in the large- and the low-recoil regions. To avoid the charmonium resonances, our choices for the large recoil region is $4m_{\mu}^2 \leq q^2 \leq 7$GeV$^2$. In the large-recoil region the relevant transversity aplitudes are (2.11)-(2.18). Since $m_{\mu}^2/q^2$ is negligible for $m_{\mu}^2 < q^2$, to simplify our numerical codes, we neglect the helicity suppressed terms of the order $O(m_{\mu}^2/q^2)$. The bands correspond to the uncertainties in the form factors only.

Before we present our numerical analysis in the low recoil, few comments are in order regarding our results. The perturbative Wilson coefficients $C_{T,9}$ do not account for the broad charmonium resonances present in the low-recoil region and are responsible for the violation of the quark-hadron duality [48]. These effects are however beyond the neglected orders in the OPE based approach. In the $B \to K \ell^+\ell^-$ decay, the duality violation in the integrated decay rate over the high-$q^2$ region is estimated to be around 2% [49]. Such an analysis for the present mode is beyond the scope of this paper and we ignore it. Since the OPE does not capture the non-perturbative effects of the local resonance structures, the actual distributions may be locally be away from the OPE predictions.

Our choice for the low-recoil phase space is $14.2 \leq q^2 \leq (m_{\Lambda_b} - m_{\Lambda^*})^2$. In the low-recoil region the masses of the muon has negligible effect and hence we put $m_{\mu} = 0$. Due to the HQET spin flavor symmetry, the amplitudes $B_{\perp\parallel}^{L(R)}$ and $B_{\parallel\perp}^{L(R)}$ vanish and the rest of the amplitudes depend on one form factors each (see Sec. 5.2). In the SM basis four form factors $f_{V,A}^{V,A}$ contribute. In addition to the uncorrelated 30% uncertainties present on the form factors, we display other dominant uncertainties for numerical illustration.
Figure 1. The differential branching ratio, lepton-side forward-backward asymmetry, and the longitudinal polarization fraction of $\Lambda_b \to \Lambda^*(\to N\bar{K})\mu^+\mu^-$ in the SM at the large- and low-recoil. The blue, green, and the red bands correspond to the uncertainties coming from the form factors, corrections to the Isgur-Wise relations, and sub-leading corrections to the amplitudes (see text for details).

The other dominant uncertainties include 5% corrections sub-leading corrections of order $O(\alpha_s\Lambda_{QCD}/m_b)$ to the transversity amplitudes in $A_{\perp,\parallel}^{L(R)}$, and 10% corrections of order $O(\Lambda_{QCD}/m_b)$ and neglected kinematical terms of order $m_{\Lambda^*}/m_{\Lambda_b}$ to the amplitudes $A_{\perp,\parallel}, A_{\perp,\parallel}$ corresponding to the improved Isgur-Wise relations. These uncertainties are
**Figure 2.** The SM predictions of $K_{1c}/K_{1cc}$ and $K_{1ss}/K_{1cc}$ at low-recoil for $\Lambda_b \to \Lambda^*(\to N\overline{K})\mu^+\mu^-$. The meaning of the bands are same as in the Fig. 1.

**Figure 3.** The differential branching ratio of $\Lambda_b \to \Lambda^*(\to N\overline{K})\mu^+\mu^-$ is shown in the SM (blue solid line) and in the NP (dashed colored) at the large- and low-recoil.

included by scaling factors. Uncertainties coming from the parameters and other sources are ignored due to the crude estimate of form factor uncertainties. In Fig. 1 we have shown
Figure 4. The ratios $K_{1c}/K_{1cc}$ and $K_{1sS}/K_{1cc}$ of $\Lambda_b \to \Lambda^*(\to N\bar{K})\mu^+\mu^-$ in the SM (blue solid line) and in the NP (dashed colored) at the low-recoil.

The SM estimates of the observables $d\mathcal{B}/dq^2$, $F_L$ and $A_{FB}^L$ for $\Lambda_b \to \Lambda^*(\to N\bar{K})\mu^+\mu^-$ at large- and low-recoil regions. In Fig. 2 we have shown the SM estimates of the two ratios $K_{1c}/K_{1cc}$ and $K_{1sS}/K_{1cc}$. In these figures, different sources of uncertainties are shown separately.

In figures 3 we have presented the NP sensitivities of the differential branching ratio of $\Lambda_b \to \Lambda^*(\to N\bar{K})\mu^+\mu^-$ in the large- and low-recoil regions. In figure 4 we have shown the NP sensitivities the ratios $K_{1c}/K_{1cc}$ and $K_{1sS}/K_{1cc}$ at low-recoil. In the NP plots, the lines correspond to the central values of all the inputs. To avoid clutter, in all the NP plots we have not shown the errors bands coming from form factors and other sources. With a future determination of the form factors in lattice QCD [50], our determinations can be improved.

7 Summary

In this paper we have studied some aspects of the $\Lambda_b \to \Lambda^*(\to N\bar{K})\ell^+\ell^-$ decay. The underlying $b \to s\ell^+\ell^-$ effective Hamiltonian is extended by including the chirality flipped counterparts of the Standard Model operator basis, and scalar and pseudo-scalar operator basis. We have presented a full angular analysis where we have also retained the masses of...
the final state leptons. The angular observables are expressed in terms of the transversity amplitudes. The four-fold distribution allows us to construct several observables that we study in the Standard Model and in model-independent new physics.

There are fourteen form factors that contribute to the decay. To reduce the uncertainties coming from the form factors, we have exploited the Heavy Quark Effective theory in the low-recoil. Working at the leading order in the $1/m_b$ and including the $\mathcal{O}(\alpha_s)$ corrections we have derived improved Isgur-Wise relations between the form factors. These relations correlate the form factors as a result of which the transversity amplitudes depend on single form factors. By means of these transversity amplitudes the short- and long-distance physics factorizes in the angular observables. The low-recoil factorization helps us construct observables from which short-distance physics can be tested with minimal form factor inputs. Alternatively, if new physics is not present, the form factor can be tested without interference from short-distance physics.

In the absence of any lattice QCD calculation, we have taken the form factors from quark model calculations. For our new physics analysis, we have used the model independent constraints on the new physics Wilson coefficients. Improved predictions in this decay will be possible in the future when the form factors are available from calculations in the lattice QCD.

Acknowledgements

The authors would like to thank Debajyoti Choudhury for fruitful discussions and Stefan Meinel for useful communications. DD acknowledges the DST, Govt. of India for the INSPIRE Faculty Fellowship (grant number IFA16-PH170). JD acknowledges the Council of Scientific and Industrial Research (CSIR), Govt. of India for JRF fellowship grant with File No. 09/045(1511)/2017-EMR-I.

A Decay kinematics

In this section we describe our kinematics. The lepton $\ell^-$ has momentum $q_{2\ell}$ and makes an angle $\theta_{\ell}$ with respect to the $+z$ axis in the dilepton rest frame (denoted as $2\ell$-RF). Therefore, $q_1^{\mu}$ read

$$q_1^{\mu}|_{2\ell-\text{RF}} = (E_\ell, -|q_{2\ell}| \sin \theta_{\ell}, 0, -|q_{2\ell}| \cos \theta_{\ell}),$$

$$q_2^{\mu}|_{2\ell-\text{RF}} = (E_\ell, |q_{2\ell}| \sin \theta_{\ell}, 0, |q_{2\ell}| \cos \theta_{\ell}),$$

where

$$q_{2\ell} = \beta_\ell \sqrt{q_{2\ell}^2}, \quad \beta_\ell = \sqrt{1 - \frac{4m_\ell^2}{q_{2\ell}^2}}.$$  (A.2)

Similarly, in the $N\bar{K}$ rest frame (denoted as $N\bar{K}$-RF), characterized by $k^2 = m_{\Lambda^*}^2$, the four-momentum $k_{1,2}$ read

$$k_1^{\mu}|_{N\bar{K}-\text{RF}} = (E_N, -|k_{N\bar{K}}| \sin \theta_{\Lambda^*} \cos \phi, -|k_{N\bar{K}}| \sin \theta_{\Lambda^*} \sin \phi, +|k_{N\bar{K}}| \cos \theta_{\Lambda^*}),$$

$$k_2^{\mu}|_{N\bar{K}-\text{RF}} = (E_K, |k_{N\bar{K}}| \sin \theta_{\Lambda^*} \cos \phi, |k_{N\bar{K}}| \sin \theta_{\Lambda^*} \sin \phi, -|k_{N\bar{K}}| \cos \theta_{\Lambda^*}).$$  (A.3)
The energies of the nucleon and the kaon at the $N\bar{K}$-rest frame are

$$E_N = \frac{k^2 + m_N^2 - m_K^2}{2\sqrt{k^2}}, \quad E_K = \frac{k^2 + m_K^2 - m_N^2}{2\sqrt{k^2}},$$

(A.4)

and the $N\bar{K}$ system four-momentum $k_{N\bar{K}}$ is

$$k_{N\bar{K}} = \frac{\sqrt{(k^2, m_N^2, m_K^2)}}{2\sqrt{k^2}}.$$  

(A.5)

### B Polarizations of the virtual gauge boson

In the dilepton rest frame, the virtual gauge boson polarization four-vectors are

$$\bar{\epsilon}^\mu(\pm) = \frac{1}{\sqrt{2}}(0, \mp 1, -i, 0), \quad \bar{\epsilon}^\mu(0) = (0, 0, 0, 1), \quad \bar{\epsilon}^\mu(t) = (1, 0, 0, 0).$$

(B.1)

The vectors satisfy the following orthonormality and completeness relations

$$\bar{\epsilon}^\mu(n)\bar{\epsilon}_\nu(n') = g_{nn'}, \quad \sum_{n,n'} \bar{\epsilon}^\mu(n)\bar{\epsilon}^\nu(n')g_{nn'} = g^{\mu\nu}, \quad n,n' = t, \pm 1, 0,$$

(B.2)

where $g_{n,n'} = \text{diag}(+1, -1, -1, -1)$ and our choice of the metric tensor is $g^{\mu\nu} = \text{diag}(1, -1, -1, -1)$.

### C Lepton helicity amplitudes

The explicit expressions of the lepton helicity amplitudes require us to calculate

$$\bar{u}_{\ell_1}(1 + \gamma_5)u_{\ell_2}, \quad \bar{\epsilon}^\mu(\lambda)\bar{\epsilon}_\mu(1 + \gamma_5)u_{\ell_2}.$$  

(C.1)

Following [51] the explicit expressions of the spinor for the lepton $\ell_1$ are

$$u_{\ell_1}(\lambda) = \begin{pmatrix} \sqrt{E_\ell + m_\ell\chi_\ell^u} \\ 2\lambda\sqrt{E_\ell - m_\ell\chi_\ell^u} \end{pmatrix}, \quad \chi_{\ell_1}^u = \begin{pmatrix} \cos \frac{\theta_\ell}{2} \\ \sin \frac{\theta_\ell}{2} \end{pmatrix},$$

(C.2)

$$\chi_{\ell_1}^{-1} = \begin{pmatrix} -\sin \frac{\theta_\ell}{2} \\ \cos \frac{\theta_\ell}{2} \end{pmatrix}.$$  

For the second lepton $\ell_2$ which is moving in the opposite direction to $\ell_1$, the two component spinor $\chi^v$ looks like

$$\chi_{-\lambda}^v = \xi_\lambda\chi_{\lambda}^u, \quad \xi_\lambda = 2\lambda e^{-2i\lambda\phi}.$$  

(C.3)

Hence we have

$$v_{\ell_2}(\lambda) = \begin{pmatrix} \sqrt{E_\ell - m_\ell\chi_{-\lambda}^v} \\ -2\lambda\sqrt{E_\ell + m_\ell\chi_{-\lambda}^v} \end{pmatrix}, \quad \chi_{\ell_2}^v = \begin{pmatrix} \sin \frac{\theta_\ell}{2} \\ -\cos \frac{\theta_\ell}{2} \end{pmatrix},$$

$$\chi_{-\ell_2}^v = \begin{pmatrix} \cos \frac{\theta_\ell}{2} \\ \sin \frac{\theta_\ell}{2} \end{pmatrix}.$$  

(C.4)
With these expressions for the spinors we obtain the following expression of the lepton helicity amplitudes

\[ L_{L}^{\frac{1}{2}+\frac{1}{2}} = -L_{R}^{-\frac{1}{2}-\frac{1}{2}} = \sqrt{q^2}(1 + \beta_L), \quad L_{L}^{-\frac{1}{2}-\frac{1}{2}} = -L_{R}^{\frac{1}{2}+\frac{1}{2}} = \sqrt{q^2}(1 - \beta_L), \]

\[ L_{L_{-1}}^{\frac{1}{2}+\frac{1}{2}} = L_{R_{-1}}^{-\frac{1}{2}-\frac{1}{2}} = L_{L_{-1}}^{-\frac{1}{2}-\frac{1}{2}} = L_{R_{+1}}^{\frac{1}{2}+\frac{1}{2}} = \sqrt{2}m_\ell \sin \theta_\ell, \]

\[ L_{L_{-1}}^{-\frac{1}{2}-\frac{1}{2}} = L_{R_{-1}}^{\frac{1}{2}+\frac{1}{2}} = -\sqrt{2}m_\ell \sin \theta_\ell, \]

\[ L_{L_{-1}}^{\frac{1}{2}+\frac{1}{2}} = L_{R_{-1}}^{-\frac{1}{2}-\frac{1}{2}} = -\sqrt{q^2}(1 - \beta_\ell)(1 - \cos \theta_\ell), \]

\[ L_{L_{-1}}^{-\frac{1}{2}-\frac{1}{2}} = L_{R_{-1}}^{\frac{1}{2}+\frac{1}{2}} = \frac{q^2}{2}(1 + \beta_\ell)(1 - \cos \theta_\ell), \]

\[ L_{R_{-1}}^{-\frac{1}{2}-\frac{1}{2}} = L_{L_{-1}}^{\frac{1}{2}+\frac{1}{2}} = \frac{q^2}{2}(1 + \beta_\ell)(1 - \cos \theta_\ell), \]

\[ L_{R_{-1}}^{-\frac{1}{2}-\frac{1}{2}} = L_{L_{-1}}^{\frac{1}{2}+\frac{1}{2}} = \sqrt{2}m_\ell \cos \theta_\ell, \]

\[ L_{L_{0}}^{-\frac{1}{2}-\frac{1}{2}} = L_{R_{0}}^{\frac{1}{2}+\frac{1}{2}} = -L_{L_{0}}^{\frac{1}{2}+\frac{1}{2}} = L_{R_{0}}^{-\frac{1}{2}-\frac{1}{2}} = 2m_\ell \cos \theta_\ell, \]

\[ L_{L_{0}}^{-\frac{1}{2}-\frac{1}{2}} = L_{R_{0}}^{\frac{1}{2}+\frac{1}{2}} = -\sqrt{q^2}(1 - \beta_\ell) \sin \theta_\ell, \]

\[ L_{L_{0}}^{\frac{1}{2}+\frac{1}{2}} = L_{R_{0}}^{-\frac{1}{2}-\frac{1}{2}} = -\sqrt{q^2}(1 + \beta_\ell) \sin \theta_\ell, \]

\[ L_{L_{t}}^{-\frac{1}{2}-\frac{1}{2}} = L_{R_{t}}^{\frac{1}{2}+\frac{1}{2}} = -L_{L_{t}}^{\frac{1}{2}+\frac{1}{2}} = 2m_\ell. \]

### D \ Λ_{b} \rightarrow \Lambda^{*} matrix elements

The \( \Lambda_{b} \rightarrow \Lambda^{*} \) transition matrix element can be parametrized using the \( \Lambda_{b} \) and \( \Lambda^{*} \) spinors in terms of form factors. For the transition through vector, and axial vector currents, the matrix elements are parametrized in terms of eight form factors [17]

\[ \langle \Lambda^{*} | \bar{s} \gamma^{\mu} b | \Lambda_{b} \rangle = \overline{u}_{\alpha}(k, s_{\Lambda^{*}}) \left\{ p^{\alpha} \left[ f^{V}_{\perp}(q^2)(m_{\Lambda_{b}} - m_{\Lambda^{*}}) \frac{q^{\mu}}{q^2} ight. ight. \\
+ f^{V}_{0}(q^2) \left. \frac{m_{\Lambda_{b}} + m_{\Lambda^{*}}}{s_{+}}(p^{\mu} + k^{\mu} - \frac{q^{\mu}}{q^2}(m_{\Lambda_{b}}^2 - m_{\Lambda^{*}}^2)) \right. \\
+ f^{A}_{\perp}(q^2) \left. \left[ (\gamma^{\mu} - 2 \frac{m_{\Lambda_{b}}}{s_{+}} p^{\mu} - 2 \frac{m_{\Lambda_{b}}}{s_{+}} k^{\mu}) \right. \right. \\
+ f^{A}_{\perp}(q^2) \left. \left. \left. \frac{g^{\mu}}{s_{+}} + m_{\Lambda^{*}} \frac{p^{\alpha}}{s_{-}} \left( \gamma^{\mu} - 2 \frac{k^{\mu}}{m_{\Lambda^{*}}} + 2 \frac{m_{\Lambda^{*}} k^{\mu} + m_{\Lambda_{b}}}{s_{-}} \right) \right. \right. \\
\right] \right\} u(p, s_{\Lambda_{b}}), \]

\[ \langle \Lambda^{*} | \bar{s} \gamma^{\mu} \gamma^{5} b | \Lambda_{b} \rangle = -\overline{u}_{\alpha}(k, s_{\Lambda^{*}}) \gamma^{5} \left\{ p^{\alpha} \left[ f^{A}_{\perp}(q^2)(m_{\Lambda_{b}} + m_{\Lambda^{*}}) \frac{q^{\mu}}{q^2} ight. ight. \\
+ f^{A}_{0}(q^2) \left. \frac{m_{\Lambda_{b}} - m_{\Lambda^{*}}}{s_{-}}(p^{\mu} + k^{\mu} - \frac{q^{\mu}}{q^2}(m_{\Lambda_{b}}^2 - m_{\Lambda^{*}}^2)) \right. \\
+ f^{A}_{\perp}(q^2) \left. \left. \left. \left[ (\gamma^{\mu} + 2 \frac{m_{\Lambda_{b}}}{s_{-}} p^{\mu} - 2 \frac{m_{\Lambda_{b}}}{s_{-}} k^{\mu}) \right. \right. \right. \\
+ f^{A}_{\perp}(q^2) \left. \left. \left. \left. \frac{g^{\mu}}{s_{+}} - m_{\Lambda^{*}} \frac{p^{\alpha}}{s_{-}} \left( \gamma^{\mu} + 2 \frac{k^{\mu}}{m_{\Lambda^{*}}} - 2 \frac{m_{\Lambda^{*}} k^{\mu} - m_{\Lambda_{b}}}{s_{-}} \right) \right. \right. \right. \\
\right] \right\} u(p, s_{\Lambda_{b}}). \]
Applying the equation of motion on (D.1) we get the matrix elements for the $\bar{s}b$ and $\bar{s}\gamma_5b$ as

$$\langle A^*|\bar{s}b|\Lambda_b\rangle = f_t^V \frac{m_{\Lambda_b} - m_{\Lambda^*}}{m_b - m_s} u_\alpha(k, s_{\Lambda^*}) p^\alpha u(p, s_{\Lambda_b}),$$  \hspace{1cm} (D.3)

$$\langle A^*|\bar{s}\gamma_5\gamma^5 b|\Lambda_b\rangle = f_t^A \frac{m_{\Lambda_b} + m_{\Lambda^*}}{m_b + m_s} u_\alpha(k, s_{\Lambda^*}) p^\alpha \gamma_5 u(p, s_{\Lambda_b}).$$  \hspace{1cm} (D.4)

The matrix elements corresponding to the tensor and axil-tensor currents are parametrized in terms of six form factors

$$\langle A^*|\bar{s}\sigma^\mu\nu q_v b|\Lambda_b\rangle = -\bar{u}_\alpha(k, s_{\Lambda^*}) \left\{ p^\alpha \left[ f_0^T(q^2) \frac{q^2}{s_+} (\gamma^\mu + k^\mu - \frac{q^\mu}{q^2} m_{\Lambda_b}^2 - m_{\Lambda^*}^2) \right] + f_1^T(q^2) (m_{\Lambda_b} + m_{\Lambda^*}) (\gamma^\mu - 2 \frac{m_{\Lambda^*}}{s_+} k^\mu) \right\} u(p, s_{\Lambda_b}),$$  \hspace{1cm} \hspace{1cm} (D.5)

$$\langle A^*|\bar{s}\sigma^\mu\nu q_5 v b|\Lambda_b\rangle = -\bar{u}_\alpha(k, s_{\Lambda^*}) \gamma_5 \left\{ p^\alpha \left[ f_0^{T5}(q^2) \frac{q^2}{s_-} (\gamma^\mu + k^\mu - \frac{q^\mu}{q^2} m_{\Lambda_b}^2 - m_{\Lambda^*}^2) \right] + f_1^{T5}(q^2) (m_{\Lambda_b} - m_{\Lambda^*}) (\gamma^\mu + 2 \frac{m_{\Lambda^*}}{s_-} k^\mu - 2 \frac{m_{\Lambda_b}^2}{s_-} k^\mu) \right\} u(p, s_{\Lambda_b}).$$  \hspace{1cm} \hspace{1cm} (D.6)

**E Helicity to Transversity**

Using the definitions of the helicity amplitudes (2.9)-(2.10), we construct the transversity amplitudes for the (axial-)vector currents

$$B_{\perp,1}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{\Lambda^+,+}^{L(R)}(-1/2, -3/2) + H_{\Lambda^+,0}^{L(R)}(+1/2, +3/2) \right],$$

$$B_{\perp,1}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{\Lambda^+,+}^{L(R)}(-1/2, -3/2) - H_{\Lambda^+,0}^{L(R)}(+1/2, +3/2) \right],$$

$$A_{\perp,0}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{\Lambda^+,0}^{L(R)}(+1/2, +1/2) + H_{\Lambda^+,0}^{L(R)}(-1/2, -1/2) \right],$$

$$A_{\perp,0}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{\Lambda^+,0}^{L(R)}(+1/2, +1/2) - H_{\Lambda^+,0}^{L(R)}(-1/2, -1/2) \right],$$

$$A_{\perp,1}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{\Lambda^+,+}^{L(R)}(+1/2, -1/2) + H_{\Lambda^+,+}^{L(R)}(-1/2, +1/2) \right],$$

$$A_{\perp,1}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{\Lambda^+,+}^{L(R)}(+1/2, -1/2) - H_{\Lambda^+,+}^{L(R)}(-1/2, +1/2) \right],$$

$$A_{\perp,1}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{\Lambda^+,+}^{L(R)}(+1/2, +1/2) + H_{\Lambda^+,+}^{L(R)}(-1/2, +1/2) \right],$$

$$A_{\perp,1}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{\Lambda^+,+}^{L(R)}(+1/2, +1/2) - H_{\Lambda^+,+}^{L(R)}(-1/2, +1/2) \right],$$

$$A_{\parallel,1}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{\Lambda^+,+}^{L(R)}(+1/2, +1/2) + H_{\Lambda^+,+}^{L(R)}(-1/2, -1/2) \right],$$

$$A_{\parallel,1}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{\Lambda^+,+}^{L(R)}(+1/2, +1/2) - H_{\Lambda^+,+}^{L(R)}(-1/2, -1/2) \right],$$

$$A_{\parallel,1}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{\Lambda^+,+}^{L(R)}(+1/2, +1/2) + H_{\Lambda^+,+}^{L(R)}(-1/2, -1/2) \right],$$

$$A_{\parallel,1}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{\Lambda^+,+}^{L(R)}(+1/2, +1/2) - H_{\Lambda^+,+}^{L(R)}(-1/2, -1/2) \right],$$

$$A_{\parallel,0}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{\Lambda^+,0}^{L(R)}(+1/2, +1/2) + H_{\Lambda^+,0}^{L(R)}(-1/2, -1/2) \right],$$

$$A_{\parallel,0}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{\Lambda^+,0}^{L(R)}(+1/2, +1/2) - H_{\Lambda^+,0}^{L(R)}(-1/2, -1/2) \right],$$

$$A_{\parallel,0}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{\Lambda^+,0}^{L(R)}(+1/2, +1/2) + H_{\Lambda^+,0}^{L(R)}(-1/2, -1/2) \right],$$

$$A_{\parallel,0}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{\Lambda^+,0}^{L(R)}(+1/2, +1/2) - H_{\Lambda^+,0}^{L(R)}(-1/2, -1/2) \right],$$
and for the scalar and pseudo-scalar operators

\[
A_{\perp S}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{SP}^{L(R)}(+1/2, +1/2) + H_{SP}^{L(R)}(-1/2, -1/2) \right].
\]

\[
A_{\parallel S}^{L(R)} = \frac{N}{\sqrt{2}} \left[ H_{SP}^{L(R)}(+1/2, +1/2) - H_{SP}^{L(R)}(-1/2, -1/2) \right].
\]

\[ \text{(E.2)} \]

F The Rarita-Schwinger spinor solutions

The solutions of Rarita-Schwinger spinors in the Λ* rest frame are [52]

\[
U_{-3/2}^1 = \sqrt{m_{\Lambda^*}} \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & -i & 0 & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix}, \quad U_{-1/2}^1 = \sqrt{\frac{m_{\Lambda^*}}{3}} \begin{pmatrix} 0 & 0 & 0 \\ 1 & 0 & 0 \\ -i & 0 & 0 \\ 0 & 2 & 0 \end{pmatrix},
\]

\[ \text{(F.1)} \]

\[
U_{+1/2}^1 = \sqrt{\frac{m_{\Lambda^*}}{3}} \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & -1 & 0 & 0 \\ 0 & -i & 0 & 0 \\ 2 & 0 & 0 & 0 \end{pmatrix}, \quad U_{+3/2}^1 = \sqrt{m_{\Lambda^*}} \begin{pmatrix} 0 & 0 & 0 \\ -1 & 0 & 0 \\ -i & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}.
\]

The solutions of the Dirac spinors corresponding to the nucleon are [51]

\[
u_{+1/2} = \frac{1}{2m_{\Lambda^*}} \begin{pmatrix} \sqrt{r_+} \cos \frac{\theta_{\Lambda^*}}{2} \\ -\sqrt{r_+} \sin \frac{\theta_{\Lambda^*}}{2} e^{i\phi} \\ \sqrt{r_-} \cos \frac{\theta_{\Lambda^*}}{2} \\ -\sqrt{r_-} \sin \frac{\theta_{\Lambda^*}}{2} e^{i\phi} \end{pmatrix}, \quad \nu_{-1/2} = \frac{1}{2m_{\Lambda^*}} \begin{pmatrix} \sqrt{r_+} \sin \frac{\theta_{\Lambda^*}}{2} e^{-i\phi} \\ \sqrt{r_+} \cos \frac{\theta_{\Lambda^*}}{2} \\ -\sqrt{r_-} \sin \frac{\theta_{\Lambda^*}}{2} e^{-i\phi} \\ -\sqrt{r_-} \sin \frac{\theta_{\Lambda^*}}{2} \end{pmatrix}.
\]

\[ \text{(F.2)} \]

G Angular Coefficients

The expression of the angular coefficients are

\[
K_{1c} = -\beta_l \left( \text{Re} (A_{-1}^L A_{1}^{L*}) - \{ L \leftrightarrow R \} \right),
\]

\[
K'_{1c} = \beta_l \left( \text{Re} (A_{-1}^L A_{0}^{L*}) + \text{Re} (A_{-1}^R A_{0}^{L*}) + \text{Re} (A_{1}^L A_{-1}^{L*}) + \text{Re} (A_{1}^R A_{-1}^{L*}) + \{ L \leftrightarrow R \} \right),
\]

\[ \text{K''}_{1c} = 0, \]

\[ \text{K}_{1cc} = \frac{1}{2} \left( |A_{-1}^L|^2 + |A_{1}^L|^2 + |A_{1}^L|^2 + |A_{-1}^L|^2 + \{ L \leftrightarrow R \} \right), \]

\[ \text{G.1} \]
\[ \mathcal{K}'_{1cc} = \left( -\text{Re}(A_{||}^R A_{||}^L)^* + \text{Re}(A_{|||S}^R A_{|||S}^L)^* - \text{Re}(A_{||}^R A_{|||S}^L)^* + \text{Re}(A_{|||S}^L A_{||}^L)^* \right) + \{L \leftrightarrow R\}, \]
\[ \mathcal{K}''_{1cc} = \left( |A_{||}^L|^2 - |A_{|||S}^L|^2 + |A_{||}^L|^2 + |A_{|||S}^L|^2 - |A_{||}^L|^2 - |A_{|||S}^L|^2 \right) + \{L \leftrightarrow R\}, \]
\[ \mathcal{K}'_{1ss} = \left( -\text{Re}(A_{||}^R A_{||}^L)^* + \text{Re}(A_{|||S}^R A_{|||S}^L)^* - \text{Re}(A_{||}^R A_{|||S}^L)^* \right) + \{L \leftrightarrow R\}, \]
\[ \mathcal{K}''_{1ss} = \left( -|A_{||}^L|^2 + |A_{||}^L|^2 + |A_{||}^L|^2 + |A_{|||S}^L|^2 \right) + \{L \leftrightarrow R\}, \]
\[ \mathcal{K}_2c = \frac{1}{4} \beta_{1c} \left( \text{Re}(A_{||}^L A_{||}^L)^* + 3\text{Re}(B_{||}^L B_{||}^R)^* \right), \]
\[ \mathcal{K}_2c = \frac{1}{4} \beta_{1c} \left( \text{Re}(A_{||}^L A_{||}^L)^* + \text{Re}(A_{||}^L A_{||}^L)^* + \text{Re}(A_{||}^L A_{||}^L)^* + \{L \leftrightarrow R\} \right), \]
\[ \mathcal{K}_2c = 0, \]
\[ \mathcal{K}_2c = \frac{1}{8} \left( |A_{||}^L|^2 + |A_{||}^L|^2 + 3|B_{||}^L|^2 + |A_{||}^L|^2 + |A_{||}^L|^2 + 3|B_{||}^L|^2 + \{L \leftrightarrow R\} \right), \]
\[ \mathcal{K}_2c = \left( -\text{Re}(A_{||}^R A_{||}^L)^* + \text{Re}(A_{||}^R A_{||}^L)^* \right) + \{L \leftrightarrow R\}, \]
\[ \mathcal{K}_2c = \left( -|A_{||}^L|^2 - |A_{||}^L|^2 - |A_{||}^L|^2 + |A_{||}^L|^2 - |A_{||}^L|^2 + |A_{||}^L|^2 \right) + \{L \leftrightarrow R\}, \]
\[ \mathcal{K}_2c = \left( 2|A_{||}^L|^2 + 2|A_{||}^L|^2 + 2|A_{||}^L|^2 + 2|A_{||}^L|^2 \right) + \{L \leftrightarrow R\}, \]
\[ \mathcal{K}_2c = \frac{1}{16} \left( 2|A_{||}^L|^2 + 2|A_{||}^L|^2 + 2|A_{||}^L|^2 + 2|A_{||}^L|^2 \right) + \{L \leftrightarrow R\}, \]
\[ \mathcal{K}_2c = \frac{1}{4} \left( \text{Re}(A_{||}^R A_{||}^L)^* - \text{Re}(A_{||}^R A_{||}^L)^* + \text{Re}(A_{||}^R A_{||}^L)^* \right) + \{L \leftrightarrow R\}, \]
\[ \mathcal{K}_2c = \frac{1}{4} \left( -|A_{||}^L|^2 - |A_{||}^L|^2 + |A_{||}^L|^2 - |A_{||}^L|^2 + |A_{||}^L|^2 \right) + \{L \leftrightarrow R\}, \]
\[ + \text{Re}(A_{10}^R A_{11}^L) + \text{Re}(A_{11}^R A_{10}^L) + 2\sqrt{3}\text{Re}(B_{11}^L A_{10}^L) - \text{Re}(A_{10}^R A_{11}^L) \\
- \text{Re}(A_{10}^R A_{11}^L) + \text{Re}(A_{11}^R A_{10}^L) + \text{Re}(A_{11}^R A_{10}^L) - 2\sqrt{3}\text{Re}(B_{11}^L A_{10}^L) - \text{Re}(A_{10}^R A_{11}^L) \\
- \text{Re}(A_{10}^R A_{11}^L) + 3\text{Re}(B_{11}^R B_{10}^L) + 3\text{Re}(B_{11}^R B_{11}^L) + \{L \leftrightarrow R\} \right), \quad (G.6) \]

\[ \mathcal{K}_{3ss} = \sqrt{3} \left( \text{Re}(B_{11}^L A_{11}^L) - \text{Re}(B_{11}^L A_{11}^L) + \{L \leftrightarrow R\} \right), \]

\[ \mathcal{K}_{3ss}' = 0, \]

\[ \mathcal{K}_{3ss}'' = -\sqrt{3} \left( \text{Re}(B_{11}^L A_{11}^L) - \text{Re}(B_{11}^L A_{11}^L) + \{L \leftrightarrow R\} \right), \quad (G.7) \]

\[ \mathcal{K}_{4ss} = \frac{\sqrt{3}}{4} \left( \text{Im}(B_{11}^L A_{11}^L) - \text{Im}(B_{11}^L A_{11}^L) + \{L \leftrightarrow R\} \right), \]

\[ \mathcal{K}_{4ss}' = 0, \]

\[ \mathcal{K}_{4ss}'' = -\sqrt{3} \left( \text{Im}(B_{11}^L A_{11}^L) - \text{Im}(B_{11}^L A_{11}^L) + \{L \leftrightarrow R\} \right), \quad (G.8) \]

\[ \mathcal{K}_{5s} = -\frac{\sqrt{6}}{4} \beta_l \left( \text{Re}(B_{11}^L A_{11}^L) - \text{Re}(B_{11}^L A_{11}^L) - \{L \leftrightarrow R\} \right), \]

\[ \mathcal{K}_{5s}' = -\frac{\sqrt{6}}{4} \beta_l \left( \text{Re}(B_{11}^L A_{11}^L) - \text{Re}(B_{11}^L A_{11}^L) + \{L \leftrightarrow R\} \right), \]

\[ \mathcal{K}_{5s}'' = 0, \quad (G.9) \]

\[ \mathcal{K}_{5sc} = \frac{\sqrt{6}}{4} \left( \text{Re}(B_{11}^L A_{11}^L) - \text{Re}(B_{11}^L A_{11}^L) + \{L \leftrightarrow R\} \right), \]

\[ \mathcal{K}_{5sc}' = 0, \quad (G.10) \]

\[ \mathcal{K}_{5sc}'' = -\sqrt{6} \left( \text{Re}(B_{11}^L A_{11}^L) - \text{Re}(B_{11}^L A_{11}^L) + \{L \leftrightarrow R\} \right), \]

\[ \mathcal{K}_{6s} = -\frac{\sqrt{6}}{4} \beta_l \left( \text{Im}(B_{11}^L A_{11}^L) - \text{Im}(B_{11}^L A_{11}^L) - \{L \leftrightarrow R\} \right), \]

\[ \mathcal{K}_{6s}' = -\frac{\sqrt{6}}{4} \beta_l \left( \text{Im}(B_{11}^L A_{11}^L) - \text{Im}(B_{11}^L A_{11}^L) + \{L \leftrightarrow R\} \right), \]

\[ \mathcal{K}_{6s}'' = 0, \quad (G.11) \]

\[ \mathcal{K}_{6sc} = \frac{\sqrt{6}}{4} \left( \text{Im}(B_{11}^L A_{11}^L) - \text{Im}(B_{11}^L A_{11}^L) + \{L \leftrightarrow R\} \right), \]

\[ \mathcal{K}_{6sc}' = 0, \]

\[ \mathcal{K}_{6sc}'' = -\sqrt{6} \left( \text{Im}(B_{11}^L A_{11}^L) - \text{Im}(B_{11}^L A_{11}^L) + \{L \leftrightarrow R\} \right). \quad (G.12) \]

We agree with the SM results presented in Ref. [17] up to the following differences: our SM results are smaller by a factor of 2 than the results presented in Ref. [17], and we differ in sign for the angular coefficients \( \mathcal{K}_{5s}, \mathcal{K}_{5sc}, \mathcal{K}_{6s}, \mathcal{K}_{6sc} \).
H Numerical inputs

The inputs used in our analysis are given in the table below.

| inputs              | values                          | inputs              | values                          |
|---------------------|---------------------------------|---------------------|---------------------------------|
| $\alpha_e(m_Z)$     | $1/127.925(16)$ [53]            | $|V_{tb}V_{ts}^*|$   | $0.0401 \pm 0.0010$ [54]        |
| $\mathcal{B}_{\Lambda}$ | $0.45 \pm 0.01$ [53]            | $m_{\Lambda_b}$    | $5.619$ GeV [53]                |
| $\mu$               | $4.8$ GeV [42]                  | $m_{\Lambda^*}$    | $1.5195$ GeV [53]               |
| $m_0(MS)$           | $4.2$ GeV [42]                  | $\tau_{\Lambda_b}$ | $(1.470 \pm 0.010) \times 10^{-12}$s [53] |
| $\alpha_s(M_Z)$     | $0.1181 \pm 0.0011$ [53]        | $\mathcal{B}_{\Lambda^*}$ | $0.45 \pm 0.01$ [53]            |

Table 1: List of inputs and their values.

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