

Unraveling new physics effects in $b \rightarrow s\ell_1\ell_2$ transitions with a model-independent perspective

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Motivated by recent anomalies in observables associated with flavor-changing neutral current transitions, specifically $b \rightarrow s\ell^+\ell^-$ processes, we present a comprehensive analysis of lepton flavor-violating (LFV) decay modes mediated by $b \rightarrow s\ell_1\ell_2$ transitions with $\ell_1 \neq \ell_2$. While such LFV processes are forbidden within the standard model, they naturally arise in several of its extensions, including models featuring additional vectorlike fermions and extra Z' bosons. Employing the most general effective Hamiltonian for $b \rightarrow s\ell_1\ell_2$ transitions, we derive the angular distributions of the relevant decay modes. Adopting a model-independent framework, we systematically study the LFV decays $B \rightarrow K^*\ell_1\ell_2$, $B_s \rightarrow \phi\ell_1\ell_2$, $B \rightarrow K_2^*\ell_1\ell_2$, and $\Lambda_b \rightarrow \Lambda\ell_1\ell_2$. Although LFV mesonic decays have been widely explored, the corresponding baryonic decays remain comparatively underinvestigated. We provide bounds on the branching ratio ($\mathcal{B}$), forward-backward asymmetry ($\mathcal{A}_{FB}$), and longitudinal lepton polarization fraction ($\mathcal{F}_L$). Furthermore, considering the projected sensitivities of the LHCb upgrade and Belle II experiments, we estimate upper limits for these observables, offering promising avenues for probing new physics in these LFV channels.

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I. INTRODUCTION

The standard model (SM) has proven remarkably effective in describing a broad range of physical phenomena, particularly at the electroweak scale. Nevertheless, it is now widely acknowledged that the SM is incomplete, as it fails to explain several critical aspects of nature. These include the origin of neutrino masses and mixing, the composition of dark matter, and the matter-antimatter asymmetry observed in the universe. To probe these open questions, flavor physics serves as a key tool. In particular, the analysis of rare decays mediated by flavor-changing neutral current (FCNC) transitions provides a promising avenue for discovering new physics (NP) beyond the SM. Significant discrepancies between experimental observations and SM predictions in these processes would signal potential contributions from NP.

Among these rare processes, the FCNC transitions $b \rightarrow s\ell\ell$ are of particular interest due to their rich phenomenology and the availability of theoretically clean observables. Ratios pertaining to lepton flavor universality (LFU), such as R_K and R_{K^*} —defined as the ratios of branching fractions $\mathcal{B}(B \rightarrow K^{(*)}\mu\mu)/\mathcal{B}(B \rightarrow K^{(*)}ee)$ —are especially useful, as hadronic uncertainties cancel to a large extent. Similarly, LFU observables in charged-current processes like R_D and R_{D^*} (in $b \rightarrow c\ell\nu$ transitions) show a deviation of around 3σ from SM predictions. Years of collaboration between theoretical and experimental physicists have significantly refined our ability to test the SM.

Earlier measurements by the LHCb collaboration reported a 3.1σ deviation in R_K within the q^2 range $[1.1, 6.0]$ GeV² [1]. Similarly, R_{K^*} measurements in the low- q^2 bins $[0.045, 1.1]$ and $[1.1, 6.0]$ GeV² showed deviations of 2.2σ and 2.5σ , respectively [2,3]. However, the situation changed notably with the updated LHCb results announced in December 2022 [4,5], which reported

$$R_K = \begin{cases} 0.994^{+0.090}_{-0.082}(\text{stat})^{+0.027}_{-0.029}(\text{syst}), & 0.045 \leq q^2 \leq 1.1 \text{ GeV}^2, \\ 0.949^{+0.042}_{-0.041}(\text{stat})^{+0.023}_{-0.023}(\text{syst}), & 1.1 \leq q^2 \leq 6.0 \text{ GeV}^2, \end{cases} \quad (1)$$

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$$R_{K^*} = \begin{cases} 0.927_{-0.087}^{+0.093} (\text{stat})_{-0.033}^{+0.034} (\text{syst}), & 0.045 \leq q^2 \leq 1.1 \text{ GeV}^2, \\ 1.027_{-0.068}^{+0.072} (\text{stat})_{-0.027}^{+0.027} (\text{syst}), & 1.1 \leq q^2 \leq 6.0 \text{ GeV}^2. \end{cases} \quad (2)$$

These updated values are more or less consistent with the SM predictions, with deviations reduced to approximately 0.2σ . Despite this, the presence of NP cannot be definitively excluded, as anomalies persist in other observables. For example, measurements of the angular observable P'_5 from LHCb [6,7] and ATLAS [8] report deviations up to 3.3σ , while CMS [9] and Belle [10] show discrepancies of about 1σ and 2.6σ , respectively.

These inconsistencies are collectively known as B anomalies, which highlight the potential involvement of NP, especially in semileptonic B meson decays [11–25]. Unlike lepton flavor-conserving decays, lepton flavor-violating (LFV) transitions offer a particularly clean probe for NP, as they are forbidden in the SM.

Various LFV processes have been studied both in the charged lepton sector (e.g., $\ell_i \rightarrow \ell_j \gamma$, $\ell_i \rightarrow \ell_j \ell_k \bar{\ell}_k$) and in B meson decays via $b \rightarrow s \ell_i \ell_j$ transitions [26–30], showing attention to radiative LFV modes as experimentally viable alternatives to purely leptonic channels, underscoring their importance in probing new physics. Experimentally, only upper bounds exist so far. LHCb has set an upper limit $\mathcal{B}(B_s \rightarrow e^\pm \mu^\mp) < 6.3 \times 10^{-9}$ at the 90% CL, and $\mathcal{B}(B \rightarrow \tau \mu) < 4.2 \times 10^{-5}$ [31]. LFV searches in semileptonic modes like $B \rightarrow K \ell_1 \ell_2$ yielded $\mathcal{B}(B \rightarrow K \mu^- e^+) < 7 \times 10^{-9}$ [32], and for vector meson modes $B \rightarrow K^* \mu e$ and $B \rightarrow \phi \mu e$, limits are 6.8×10^{-9} and 10.1×10^{-9} , respectively. Despite the experimental complexity introduced by τ leptons in the final state, Belle has reported $\mathcal{B}(B^+ \rightarrow K^+ \mu^\pm \tau^\mp) < 3.9 \times 10^{-5}$ [33]. LHCb searches for $B \rightarrow K^* \mu^\pm \tau^\mp$ yielded $\mathcal{B}(B \rightarrow K^* \tau^+ \mu^-) < 1.0 \times 10^{-5}$ and $\mathcal{B}(B \rightarrow K^* \tau^- \mu^+) < 8.2 \times 10^{-6}$ [34].

B anomalies fall into two major categories: (i) deviations in τ vs. light-lepton universality in $b \rightarrow c \ell \nu$ transitions, and (ii) deviations in FCNC-mediated $b \rightarrow s \ell^+ \ell^-$ processes. These suggest the existence of NP particles and have prompted extensive theoretical and experimental exploration. Proposed NP models range from heavy vector and scalar mediators to leptoquarks and nonuniversal Z' bosons [35–65]. Common to these models is the prediction that violations of LFU may also imply LFV in B and τ decays. Currently, experimental upper bounds on LFV modes place stringent constraints on these NP scenarios. Future data will be crucial in confirming or refuting these possibilities.

To investigate LFV in $b \rightarrow s \ell_1 \ell_2$ transitions, both mesonic and baryonic decays must be considered. Despite sharing the same underlying partonic process, they offer complementary perspectives on NP. For instance, while short-distance NP effects are absent in $\Lambda_b \rightarrow \Lambda \mu^+ \mu^-$ decays [66,67], such effects appear to influence angular

observables in $B \rightarrow K^* \mu^+ \mu^-$ decays [68–71]. This suggests that decays like $\Lambda_b \rightarrow \Lambda \ell_1 \ell_2$ can provide additional insights alongside mesonic channels, such as $B^+ \rightarrow K^+ \ell_1 \ell_2$ and $\bar{B}_s \rightarrow \ell_1 \ell_2$. Moreover, the spin structure of the baryonic decays leads to a richer set of hadronic matrix elements.

Although so far there is no direct evidence of LFV, many theoretical models involve the influence of FCNC mediated by the Z boson [72,73], such as nonuniversal Z' models [74,75], leptoquark models [76], and the minimal supersymmetric standard model [77,78]. Model-independent studies of LFV in B decays involving τ leptons were initially motivated in [79,80], which established a connection between LFU violation and LFV branching ratios, while various NP scenarios [81,82] have since been proposed to account for these effects. This work focuses on LFV decays mediated by $b \rightarrow s \ell_1 \ell_2$ transitions and evaluates $\mathcal{B}$, $\mathcal{A}_{FB}$, and $\mathcal{F}_L$ in the channels $B \rightarrow K^* \ell_1 \ell_2$, $B \rightarrow \phi \ell_1 \ell_2$, $B \rightarrow K_2^* \ell_1 \ell_2$, and $\Lambda_b \rightarrow \Lambda \ell_1 \ell_2$. Constraints on NP couplings are derived from current experimental bounds $\mathcal{B}(B^+ \rightarrow K^+ \tau^\mp \mu^\pm)$ and $\mathcal{B}(B_s \rightarrow \tau^\mp \mu^\pm)$. The remainder of this paper is structured as follows. Section II introduces the effective Hamiltonian and theoretical framework for $b \rightarrow s \ell_1 \ell_2$ transitions. Section III presents the numerical analysis for $B \rightarrow K^* \ell_1 \ell_2$ and $B \rightarrow \phi \ell_1 \ell_2$ decays. Section IV focuses on $B \rightarrow K_2^* \ell_1 \ell_2$, and Sec. V discusses results for $\Lambda_b \rightarrow \Lambda \ell_1 \ell_2$. A summary of key findings is provided in Sec. VI.

II. EFFECTIVE HAMILTONIAN AND THEORETICAL ASPECTS

We commence by employing the effective Hamiltonian to describe the LFV in $b \rightarrow s \ell_1^+ \ell_2^-$ transition,

$$\mathcal{H}^{\text{eff}} = -\frac{G_F \alpha_{\text{em}}}{\sqrt{2}\pi} V_{ts}^* V_{tb} \sum_i (C_i \mathcal{O}_i + C'_i \mathcal{O}'_i), \quad (3)$$

where G_F is the Fermi constant, α_{em} is the fine-structure constant, V_{ts} and V_{tb} are the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements, and $i = S, P, V$, and A correspond to scalar, pseudoscalar, vector, and axial-vector operators, which read as

$$\begin{aligned} \mathcal{O}_S^{(\prime)} &= [\bar{s} P_{R(L)} b] [\bar{\ell}_2 \ell_1], & \mathcal{O}_P^{(\prime)} &= [\bar{s} P_{R(L)} b] [\bar{\ell}_2 \gamma_5 \ell_1] \\ \mathcal{O}_V^{(\prime)} &= [\bar{s} \gamma^\mu P_{L(R)} b] [\bar{\ell}_2 \gamma_\mu \ell_1], & \mathcal{O}_A^{(\prime)} &= [\bar{s} \gamma^\mu P_{L(R)} b] [\bar{\ell}_2 \gamma_\mu \gamma_5 \ell_1], \end{aligned} \quad (4)$$

where $P_{L,R} = (1 \mp \gamma_5)/2$ are the chiral projectors. It is important to highlight that the electromagnetic penguin operator $\mathcal{O}_7$, which is involved in $b \rightarrow s\ell\ell$ decays, is characterized by

$$\mathcal{O}_7 = \frac{e}{16\pi^2} m_b (\bar{s} \sigma_{\mu\nu} P_R b) F^{\mu\nu}. \quad (5)$$

However, it is crucial to note that this operator cannot induce lepton flavor-violating contributions because of the universal nature of electromagnetic interactions. Tensor operators such as $\mathcal{O}_T = [\bar{s} \sigma^{\mu\nu} b][\bar{\ell}_2 \sigma_{\mu\nu} \ell_1]$ are not included in this study. These are typically suppressed or absent in many UV-complete models (e.g., Z' , leptoquarks, etc.) and do not contribute to purely leptonic LFV decays like $B_s \rightarrow \ell_1 \ell_2$, which provide the strongest constraints. Their inclusion would also introduce significant hadronic uncertainties. We, therefore, focus on scalar, pseudoscalar, vector, and axial-vector operators. The coefficients $C_{S,P,V,A}^{(\prime)}$ represent short-distance Wilson coefficients that are absent in the SM but can take nonzero values in various scenarios outside the SM. In the SM, ℓ_1 and ℓ_2 are typically leptons of the same flavor, often denoted as ℓ , and the corresponding operators $\mathcal{O}_{V,A}$ are conventionally labeled as $\mathcal{O}_{9,10}$, with their respective Wilson coefficients denoted as $C_{9,10}$.

TABLE I. Experimental upper limits for LFV B decays at the 90% CL.

Observable	Upper bound
$\mathcal{B}(B^+ \rightarrow K^+ \tau^+ \mu^-)$	3.9×10^{-5} [90]
$\mathcal{B}(B^+ \rightarrow K^+ \tau^- \mu^+)$	4.5×10^{-5} [91]
$\mathcal{B}(\bar{B}_s \rightarrow \tau^\mp \mu^\pm)$	3.5×10^{-5} [92]

A. Constraints on new couplings

First, we consider the constraints on several combinations of Wilson coefficients from measurements of mesonic LFV decays. The quark-level transition for the mesonic $\bar{B}_s \rightarrow \ell_1^- \ell_2^+$ and $B \rightarrow K \ell_1^- \ell_2^+$ decays is the same as the processes we are looking for. So, we can use the upper limit on the combinations of Wilson coefficients extracted from these decay modes and can obtain the upper limits on various observables of $B_s \rightarrow \{K^*, \phi\} \ell_1 \ell_2$, $B \rightarrow K_2^* \ell_1 \ell_2$ and $\Lambda_b \rightarrow \Lambda \ell_1 \ell_2$ processes. We consider the branching ratios of the decay modes $\bar{B}_s \rightarrow \ell_1^- \ell_2^+$ and $B \rightarrow K \ell_1^- \ell_2^+$, for which the experimental upper limits at the 90% CL are reported in Table I. Utilizing Eq. (3), the branching ratios for the corresponding decay channels can be written as follows:

$$\begin{aligned} \mathcal{B}(B^+ \rightarrow K^+ \ell_1^- \ell_2^+) &= 10^{-8} \{ c_{\ell_1 \ell_2}^P |C_{P_+}|^2 + c_{\ell_1 \ell_2}^{10+} |C_{10_+}|^2 + c_{\ell_1 \ell_2}^S |C_{S_+}|^2 \\ &\quad + c_{\ell_1 \ell_2}^{9+} |C_{9_+}|^2 + c_{\ell_1 \ell_2}^{P10} \text{Re}[C_{P_+}^* C_{10_+}] + c_{\ell_1 \ell_2}^{S9} \text{Re}[C_{S_+}^* C_{9_+}] \}, \end{aligned} \quad (6)$$

and

$$\begin{aligned} \mathcal{B}(\bar{B}_s \rightarrow \ell_1^- \ell_2^+) &= \frac{\alpha_{\text{em}}^2 G_F^2 |V_{tb} V_{ts}^*|^2}{m_{B_s}^3} \frac{\tau_{B_s}}{64\pi^3} f_{B_s}^2 \lambda^{1/2}(m_{B_s}^2, m_{\ell_1}^2, m_{\ell_2}^2) \\ &\quad \times \left\{ [m_{B_s}^2 - (m_{\ell_1} - m_{\ell_2})^2] \left| \frac{m_{B_s}^2}{m_b + m_s} C_{P_-} + (m_{\ell_1} + m_{\ell_2}) C_{10_-} \right|^2 \right. \\ &\quad \left. + [m_{B_s}^2 - (m_{\ell_1} + m_{\ell_2})^2] \left| \frac{m_{B_s}^2}{m_b + m_s} C_{S_-} + (m_{\ell_1} - m_{\ell_2}) C_{9_-} \right|^2 \right\}. \end{aligned} \quad (7)$$

Both Eqs. (6) and (7) agree with previous results in the literature [83,84]. Here, we have adopted the notation $C_{Y\pm} = C_Y \pm C'_Y$. The best-fit values are obtained under the assumption $C'_Y = 0$, corresponding to purely left-handed contributions. C'_Y are not favored by current experimental data, but may become relevant in future analyses as more precise angular or baryonic observables become available. By using the mass values from the Particle Data Group (PDG) [85], the CKM factors obtained from UT-fit collaboration [86], the decay constant $f_{B_s} = 215$ MeV, as provided in [87], and relying on lattice QCD/light cone sum rule results from Refs. [88,89], we have determined the coefficients $c_{\ell_1 \ell_2}^i$ in Eq. (6), as

displayed in Table II. It is worth mentioning that these coefficients are not influenced by the charges of the leptons, with the exception of $c_{\ell_1 \ell_2}^{S9}$. It is important to note that the coefficient $c_{\ell_1 \ell_2}^{S9}$ is directly proportional to the difference in masses between ℓ_1 and ℓ_2 . As a result, this coefficient

TABLE II. The predictions for the coefficients governing $B^+ \rightarrow K^+ \ell_1^- \ell_2^+$ decays ($\ell_1^- = \mu^-$, $\ell_2^+ = \tau^+$) based on the hadronic form factors from Refs. [88,89].

$c_{\ell_1 \ell_2}^{9+}$	$c_{\ell_1 \ell_2}^{10+}$	$c_{\ell_1 \ell_2}^S$	$c_{\ell_1 \ell_2}^P$	$c_{\ell_1 \ell_2}^{S9}$	$c_{\ell_1 \ell_2}^{P10}$
1.09	1.14	1.47	1.58	-1.35	1.66

changes its sign based on the charge of the heavier lepton in the process. We should emphasize that the values presented in Table II are significantly influenced by the choice of the electromagnetic coupling constant α_{em} . In this analysis, we have used $\alpha_{\text{em}} = 1/133$. However, it is worth noting that a different value for α_{em} can be applied by appropriately scaling the $c_{\ell_1 \ell_2}^i$ coefficients.

To determine the NP parameter space consistent with current experimental data, we perform a χ^2 analysis. The χ^2 function is defined as

$$\chi^2(C_{\text{NP}}) = \sum_i \frac{(\mathcal{O}_i^{\text{Th}}(C_{\text{NP}}) - \mathcal{O}_i^{\text{Exp}})^2}{(\Delta \mathcal{O}_i^{\text{Exp}})^2 + (\Delta \mathcal{O}_i^{\text{SM}})^2}, \quad (8)$$

where $\mathcal{O}_i^{\text{Th}}$ and $\mathcal{O}_i^{\text{Exp}}$ denote the theoretical prediction and the experimental central value of the i th observable, respectively. The denominator accounts for both experimental and SM theoretical uncertainties. We employ the experimental upper limits provided in Table I and Eqs. (6) and (7) to constrain various combinations of NP couplings. It is important to note that for the decay modes $B^+ \rightarrow K^+ \mu^\pm \tau^\mp$ and $B_s \rightarrow \mu^\pm \tau^\mp$, only 90% CL upper bounds on the branching ratios are available. To incorporate these into the χ^2 analysis, we model the corresponding observable as having a central value of zero and an uncertainty given by U.L./1.645, where U.L. denotes the upper limit. Since these processes are highly suppressed in the standard model, the associated theoretical uncertainty from the SM is taken to be negligible. The NP parameter space is then obtained by minimizing the χ^2 function. The analysis is carried out using the best-fit values of the Wilson coefficients, given by $C_9 = -55.361$, $C_{10} = 2.3 \times 10^{-2}$, $C_S = -8.941$, and $C_P = 6.763$. In this work, we restrict our analysis to real-valued LFV Wilson coefficients, i.e., we neglect potential complex phases.

It is important to reiterate that we do not include $\tau\ell$ decays in this analysis, as the constraints derived from these decays resemble those from the $\tau\mu$ channel. Additionally, for the sake of simplicity, we do not account for the $\mathcal{O}_i^{\ell_1 \ell_2}$

operators. This choice is made due to the fact that these operators are less favorable when attempting to fit data from $b \rightarrow s \ell \ell$ transitions [68–71,93]. Nevertheless, it is worth emphasizing that baryonic channels exhibit different dependencies on the primed operators compared to mesonic channels. This distinction could become relevant as scenarios involving these operators gain prominence in addressing the B anomalies.

We explore two distinct two-dimensional scenarios, each of which allows only specific combinations of NP Wilson coefficients to have nonzero values. These scenarios encompass, like the first scenario, $C_9^{\ell_1 \ell_2}$ and $C_{10}^{\ell_1 \ell_2}$, with the constraint that only these two coefficients are allowed to be nonzero. The second scenario is inspired by the standard model effective field theory, where we set $C_9^{\ell_1 \ell_2} = -C_{10}^{\ell_1 \ell_2}$ and $C_S^{\ell_1 \ell_2} = -C_P^{\ell_1 \ell_2}$ as the specific constraints on the coefficients. Scenarios with simultaneous scalar and vector operators can exhibit interference effects $\propto \text{Re}(C_{S,P}^* C_{9,10})$, leading to qualitatively different phenomenology compared to single-operator cases. These effects can alter branching ratios and angular observables, and we acknowledge their relevance as a promising direction for future work.

Scenario I: $C_9^{\ell_1 \ell_2} \neq 0$, $C_{10}^{\ell_1 \ell_2} \neq 0$, $C_S^{\ell_1 \ell_2} = C_P^{\ell_1 \ell_2} = 0$.

Scenario II: $C_9^{\ell_1 \ell_2} = -C_{10}^{\ell_1 \ell_2}$, $C_S^{\ell_1 \ell_2} = -C_P^{\ell_1 \ell_2}$.

The obtained bounds for $\tau^+ \mu^-$ and $\mu^+ \tau^-$ final states are given in Fig. 1. The left two plots of the figure show C_9 versus C_{10} , while setting all others to zero, which are independent of the charge configuration in the final state, and we only consider the strongest bound of $\bar{B}_s \rightarrow \ell^+ \mu^-$ and $B \rightarrow K \tau^+ \mu^-$ decay processes. Since the interference between $C_9^{\ell_1 \ell_2}$ and $C_S^{\ell_1 \ell_2}$ relies on the charge configuration of the leptons in the final state, we have generated plots in the bottom panel for both the $\tau^+ \mu^-$ and $\tau^- \mu^+$ final states. Upon comparing the plots in Fig. 1, we observe substantial disparities between the $\tau^+ \mu^-$ and $\tau^- \mu^+$. Consequently, we emphasize the importance of conducting separate analyses for these distinct final states.

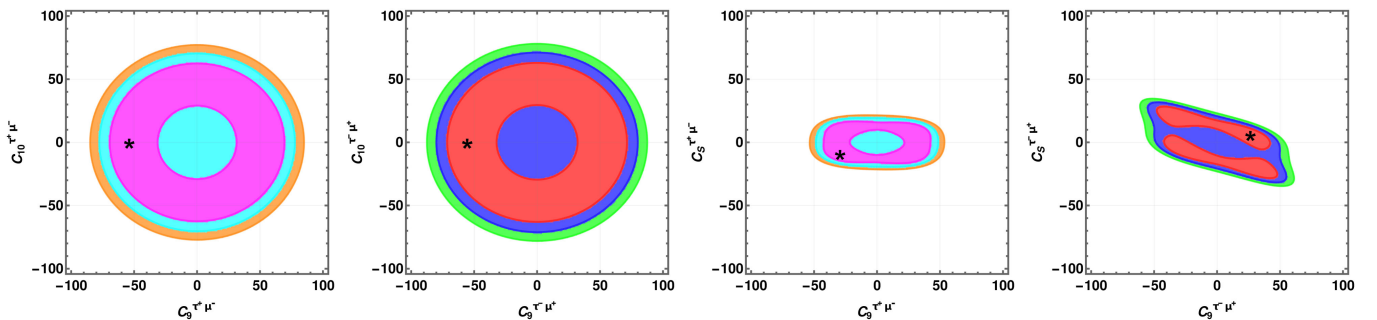


FIG. 1. The allowed parameter space is determined by combining the constraints on various Wilson coefficients derived from the 90% confidence level upper limits of the $\bar{B}_s \rightarrow \ell^+ \ell_2^+$ and $B \rightarrow K \ell_1^- \ell_2^+$ decays. In the $\tau^+ \mu^-$ channel, the magenta, cyan, and orange regions represent the 1σ , 2σ , and 3σ error bands, respectively, while for the $\tau^- \mu^+$ mode, the corresponding bands are shown in red, blue, and green.

III. DECAY DISTRIBUTION OF $B \rightarrow K^*\ell_1\ell_2$ AND $B \rightarrow \phi\ell_1\ell_2$

These processes occur through two different decay channels: $B \rightarrow K^*(\rightarrow K\pi)\ell_1\ell_2$ and $B_s \rightarrow \phi(\rightarrow K\bar{K})\ell_1\ell_2$. Since the angular distribution expression for the latter decay can be derived by straightforward substitutions of the corresponding mass and form factors in the expression of the former decay mode, our primary focus will be on $\bar{B} \rightarrow \bar{K}^*(\rightarrow K^-\pi^+)\ell_1^-\ell_2^+$. Here, we adopt the kinematics conventions, which are fine-tuned to align with the standards

adhered to in experiments conducted at the LHC [17]. Additional details regarding the kinematics of this process are provided in the Appendix A of this paper. In addition to θ_ℓ , we also consider the parameter θ_K , which signifies the angle between the decay axis oriented along the $-z$ direction and the trajectory of the K^- particle in the rest frame of $\bar{K}^*$. This is illustrated in Fig. 10. Furthermore, we use ϕ to denote the angle between the planes defined by the $K\pi$ and $\ell_1^-\ell_2^+$ systems.

In this scenario, the hadronic matrix elements involve a more extensive set of form factors, including

$$\begin{aligned} \langle \bar{K}^*(k) | \bar{s}\gamma^\mu(1-\gamma_5)b | \bar{B}(p) \rangle &= \varepsilon_{\mu\nu\rho\sigma}\varepsilon^{*\nu}p^\rho k^\sigma \frac{2V(q^2)}{m_B + m_{K^*}} - i \left[\varepsilon_\mu^*(m_B + m_{K^*})A_1(q^2) \right. \\ &\quad \left. - i(p+k)_\mu(\varepsilon^* \cdot q) \frac{A_2(q^2)}{m_B + m_{K^*}} - q_\mu(\varepsilon^* \cdot q) \frac{2m_{K^*}}{q^2} [A_3(q^2) - A_0(q^2)] \right], \\ \langle \bar{K}^*(k) | \bar{s}\sigma_{\mu\nu}q^\nu(1-\gamma_5)b | \bar{B}(p) \rangle &= 2i\varepsilon_{\mu\nu\rho\sigma}\varepsilon^{*\nu}p^\rho k^\sigma T_1(q^2) - [(\varepsilon^* \cdot q)(2p-q)_\mu - \varepsilon_\mu^*(m_B^2 - m_{K^*}^2)]T_2(q^2) \\ &\quad - (\varepsilon^* \cdot q) \left[\frac{q^2}{m_B^2 - m_{K^*}^2} (p+k)_\mu - q_\mu \right] T_3(q^2), \end{aligned} \quad (9)$$

where ε_μ is the polarization vector of K^* , and $A_{1,2,3}(q^2)$ are the form factors having the relation $2m_{K^*}A_3(q^2) = (m_B + m_{K^*})A_1(q^2) - (m_B - m_{K^*})A_2(q^2)$. The full angular distribution of the above decay processes is¹

$$\frac{d^4\mathcal{B}(B \rightarrow \bar{K}^* \rightarrow (K\pi)\ell_1^-\ell_2^+)}{dq^2 d\cos\theta_\ell d\cos\theta_K d\phi} = \frac{9}{32\pi} I(q^2, \theta_\ell, \theta_K, \phi), \quad (10)$$

with

$$\begin{aligned} I(q^2, \theta_\ell, \theta_K, \phi) &= I_1^c(q^2) + [I_1^s(q^2) - I_1^c(q^2)]\sin^2\theta_K + [I_2^s(q^2) + [I_2^s(q^2) - I_2^c(q^2)]\sin^2\theta_K] \cos 2\theta_\ell \\ &\quad + I_3(q^2)\sin^2\theta_K\sin^2\theta_\ell \cos 2\phi + I_4(q^2) \sin 2\theta_K \sin 2\theta_\ell \cos \phi \\ &\quad + I_5(q^2) \sin 2\theta_K \sin \theta_\ell \cos \phi + [I_6^c(q^2) + [I_6^s(q^2) - I_6^c(q^2)]\sin^2\theta_K] \cos \theta_\ell \\ &\quad + I_7(q^2) \sin 2\theta_K \sin \theta_\ell \sin \phi + I_8(q^2) \sin 2\theta_K \sin 2\theta_\ell \sin \phi \\ &\quad + I_9(q^2)\sin^2\theta_K\sin^2\theta_\ell \sin 2\phi. \end{aligned} \quad (11)$$

After integrating over angles, the differential decay rate is simply

$$\frac{d\mathcal{B}}{dq^2} = \frac{1}{4} [6I_1^s(q^2) + 3I_1^c(q^2) - 2I_2^s(q^2) - I_2^c(q^2)]. \quad (12)$$

Similarly, the forward-backward asymmetry and the longitudinal polarization fraction of the lepton in the decay process are defined as follows:

¹Please notice that the convention used in Eq. (9) is such that $\varepsilon_{0123} = +1$.

$$A_{FB}(q^2) = \frac{3I_6^s(q^2) + \frac{3}{2}I_6^c(q^2)}{6I_1^s(q^2) + 3I_1^c(q^2) - 2I_2^s(q^2) - I_2^c(q^2)}, \quad (13)$$

$$\mathcal{F}_L(q^2) = \frac{3I_1^c(q^2) - I_2^c(q^2)}{6I_1^s(q^2) + 3I_1^c(q^2) - 2I_2^s(q^2) - I_2^c(q^2)}. \quad (14)$$

The angular coefficients, which depend on the momentum transfer q^2 , are constructed from combinations of the decay's helicity amplitudes. These amplitudes can also be represented using the transversity amplitudes $A_{\perp,\parallel,0,t}^{L(R)} \equiv A_{\perp,\parallel,0,t}^{L(R)}(q^2)$ in the following manner:

$$\begin{aligned}
A_{\perp}^{L(R)} &= \mathcal{N}_{K^*} \sqrt{2} \lambda_B^{1/2} \left[(C_9 \mp C_{10}) \frac{V(q^2)}{m_B + m_{K^*}} \right], \\
A_{\parallel}^{L(R)} &= -\mathcal{N}_{K^*} \sqrt{2} (m_B^2 - m_{K^*}^2) \left[(C_9 \mp C_{10}) \frac{A_1(q^2)}{m_B - m_{K^*}} \right], \\
A_0^{L(R)} &= -\frac{\mathcal{N}_{K^*}}{2m_{K^*} \sqrt{q^2}} (C_9 \mp C_{10}) \left[(m_B^2 - m_{K^*}^2 - q^2)(m_B + m_{K^*}) A_1(q^2) - \frac{\lambda_B A_2(q^2)}{m_B + m_{K^*}} \right], \\
A_t^{L(R)} &= -\mathcal{N}_{K^*} \frac{\lambda_B^{1/2}}{\sqrt{q^2}} \left[(C_9 \mp C_{10}) + \frac{q^2}{m_b + m_s} \left(\frac{C_S}{m_1 - m_2} \mp \frac{C_P}{m_1 + m_2} \right) \right] A_0(q^2),
\end{aligned} \tag{15}$$

where for shortness, $\lambda_B = \lambda(m_B, m_{K^*}, \sqrt{q^2})$, $\lambda_q = \lambda(m_1, m_2, \sqrt{q^2})$, and

$$\mathcal{N}_{K^*} = V_{tb} V_{ts}^* \left[\frac{\tau_{B_d} G_F^2 \alpha^2}{3 \times 2^{10} \pi^5 m_B^3} \lambda_B^{1/2} \lambda_q^{1/2} \right]^{1/2}. \tag{16}$$

The upper signs in the aforementioned equations pertain to A_t^L , while the lower signs relate to A_t^R . It is noteworthy that A_t also carries the superscript $L(R)$, indicating the chirality of the lepton pair. This may appear unusual when juxtaposed with the lepton flavor-conserving case, and we will now provide an explanation for this. When $\ell_1 = \ell_2$, the pseudoscalar density can be reformulated as follows:

$$\bar{\ell} \gamma_5 \ell = \frac{q^\mu}{2m_\ell} (\bar{\ell} \gamma_\mu \gamma_5 \ell). \tag{17}$$

This allows us to assimilate the contributions originating from the operator $\mathcal{O}_P^{(\prime)}$ into the amplitude A_t . This amplitude is associated with the timelike polarization vector of the virtual vector boson, denoted as $\epsilon_V^\mu(t) = q^\mu / \sqrt{q^2}$. An analogous method is not applicable to the scalar operator $\mathcal{O}_S^{(\prime)}$ due to the fact that $q^\mu (\bar{\ell} \gamma_\mu \ell) = 0$. In this case, it is necessary to introduce a new amplitude, denoted as A_S , to account for the remaining scalar contribution. In LFV scenarios, where $m_1 \neq m_2$, it becomes possible to utilize the Ward identities to absorb both the scalar and pseudo-scalar densities into the vector and axial currents, respectively. Consequently, in LFV cases, the amplitudes A_t and A_S are substituted by $A_t^{L(R)}$. Although the expressions for $A_t^{L,R}$ become problematic in the limit where $m_1 = m_2$, we have verified that the angular coefficients remain well defined, and the standard formulas presented in Ref. [94] are recovered. Finally, the angular coefficients $I_{1-9}(q^2)$ in terms of the transversity amplitudes (15) are given by

$$I_1^s(q^2) = \frac{4m_1 m_2}{q^2} \text{Re}(A_{\parallel}^L A_{\parallel}^{R*} + A_{\perp}^L A_{\perp}^{R*}) + [|A_{\perp}^L|^2 + |A_{\parallel}^L|^2 + (L \rightarrow R)] \frac{\lambda_q + 2[q^4 - (m_1^2 - m_2^2)^2]}{4q^4}, \tag{18}$$

$$I_1^c(q^2) = \frac{8m_1 m_2}{q^2} \text{Re}(A_0^L A_0^{R*} - A_t^L A_t^{R*}) + [|A_0^L|^2 + |A_0^R|^2] \frac{q^4 - (m_1^2 - m_2^2)^2}{q^4} - 2 \frac{(m_1^2 - m_2^2)^2 - q^2(m_1^2 + m_2^2)}{q^4} (|A_t^L|^2 + |A_t^R|^2), \tag{19}$$

$$I_2^s(q^2) = \frac{\lambda_q}{4q^4} [|A_{\parallel}^L|^2 + |A_{\perp}^L|^2 + (L \rightarrow R)], \tag{20}$$

$$I_2^c(q^2) = -\frac{\lambda_q}{q^4} (|A_0^R|^2 + |A_0^L|^2), \tag{21}$$

$$I_3^c(q^2) = \frac{\lambda_q}{2q^4} [|A_{\perp}^L|^2 - |A_{\parallel}^L|^2 + (L \rightarrow R)], \tag{22}$$

$$I_4(q^2) = -\frac{\lambda_q}{\sqrt{2}q^4} [\text{Re}(A_{\parallel}^L A_0^{L*} + (L \rightarrow R))], \tag{23}$$

$$I_5(q^2) = \frac{\sqrt{2}\lambda_q^{1/2}}{q^2} \left[\text{Re}(A_0^L A_\perp^{L*} - (L \rightarrow R)) - \frac{m_1^2 - m_2^2}{q^2} \text{Re}(A_\parallel^L A_\parallel^{L*} + (L \rightarrow R)) \right], \quad (24)$$

$$I_6^s(q^2) = -\frac{2\lambda_q^{1/2}}{q^2} [\text{Re}(A_\parallel^L A_\perp^{L*} - (L \rightarrow R))], \quad (25)$$

$$I_6^c(q^2) = -\frac{4\lambda_q^{1/2}}{q^2} \frac{m_1^2 - m_2^2}{q^2} \text{Re}(A_0^L A_\parallel^{L*} + (L \rightarrow R)), \quad (26)$$

$$I_7(q^2) = -\frac{\sqrt{2}\lambda_q^{1/2}}{q^2} \left[\frac{m_1^2 - m_2^2}{q^2} \text{Im}(A_\perp^L A_\parallel^{L*} + (L \rightarrow R)) + \text{Im}(A_0^L A_\parallel^{L*} - (L \rightarrow R)) \right], \quad (27)$$

$$I_8(q^2) = \frac{\lambda_q}{\sqrt{2}q^4} \text{Im}(A_0^L A_\perp^{L*} + (L \rightarrow R)), \quad (28)$$

$$I_9(q^2) = -\frac{\lambda_q}{q^4} \text{Im}(A_\perp^R A_\parallel^{R*} + A_\perp^L A_\parallel^{L*}). \quad (29)$$

The form factors associated with the transversity amplitude are obtained using the light cone sum rule (LCSR) method [95] and are expressed as follows:

$$F_y(q^2) = \frac{1}{1 - q^2/m_{R,y}^2} \sum_{k=0,1,2} \beta_k [z(q^2) - z(0)]^k, \quad (30)$$

where $z(q^2) = \frac{\sqrt{t_+ - q^2} - \sqrt{t_+ - t_0}}{\sqrt{t_+ - q^2} + \sqrt{t_+ - t_0}}$, $t_\pm = (m_B \pm m_{K^*})^2$, $t_0 = t_+(1 - \sqrt{1 - t_-/t_+})$, and m_{R,F_y} is the resonance mass corresponding to the form factor. In our calculations, we use the resonance masses as follows:

$$m_{R,(V,T_1)} = 5.412 \text{ GeV}, \quad m_{R,A_0} = 5.336 \text{ GeV}, \\ m_{R,(T_2,T_3,A_1,A_3)} = 5.829 \text{ GeV}. \quad (31)$$

We apply the above formalism originally developed for $B \rightarrow K^*$ decay to analyze the process $B \rightarrow \phi\ell_1\ell_2$. This is achieved by straightforwardly substituting the relevant mass and form factors for the corresponding vector meson ϕ . The numerical values for the parameters associated with the form factors with 1σ uncertainty are given in Table III.

A. Numerical analysis for $B \rightarrow K^*(\rightarrow K\pi)\ell_1\ell_2$ and $B_s \rightarrow \phi(\rightarrow K\bar{K})\ell_1\ell_2$

To quantitatively highlight the significance of the terms that scale the Wilson coefficients, we utilize the form factors provided in Ref. [96], and distinguish between scenarios where LFV originates from vector-type and scalar-type operators. The scaling factors (multipliers) for the Wilson coefficients are obtained by integrating over the full kinematic range of the dilepton invariant mass

squared, q^2 . It is noteworthy that the integrand functions exhibit distinct behaviors: contributions from terms proportional to $|C_{9,10}|^2$ dominate in the intermediate q^2 region, while those involving $|C_{S,P}|^2$ become more significant at higher values of q^2 . To illustrate this behavior, we analyze the decay channels $\tau^+\mu^-$ and $\tau^-\mu^+$ separately.

In Fig. 2, we present the differential distributions for the $B \rightarrow K^*\tau^+\mu^-$ decay. The plots display the $\mathcal{B}$ (left), $\mathcal{A}_{FB}$ (middle), and $\mathcal{F}_L$ (right), highlighting the differences between Scenario I (green) and Scenario II (red), each shown with their corresponding 1σ uncertainty bands. Similarly, Fig. 3 shows the corresponding observables for the $B \rightarrow K^*\tau^-\mu^+$ decay, with Scenario I illustrated in magenta and Scenario II in cyan blue, again including 1σ error bands. In Fig. 4, we present the differential distributions for the $B \rightarrow \phi\tau^+\mu^-$ decay. The plots display the $\mathcal{B}$ (left), $\mathcal{A}_{FB}$ (middle), and $\mathcal{F}_L$ (right), highlighting the differences between Scenario I (purple) and

TABLE III. Form factors for $B \rightarrow K^*\ell_1\ell_2$ (top panel) and $B_s \rightarrow \phi\ell_1\ell_2$ (bottom panel).

F_i	β_0	β_1	β_2
V	0.34 ± 0.04	-1.05 ± 0.24	2.37 ± 1.39
T_1	0.28 ± 0.03	-0.89 ± 0.19	1.95 ± 1.10
T_2	0.28 ± 0.03	0.40 ± 0.18	0.36 ± 0.51
T_3	0.67 ± 0.08	1.48 ± 0.49	1.92 ± 1.96
A_0	0.36 ± 0.05	-1.04 ± 0.27	1.12 ± 1.35
A_1	0.27 ± 0.03	0.39 ± 0.19	-0.11 ± 0.48
A_3	0.26 ± 0.03	0.60 ± 0.20	0.12 ± 0.84
F_i	β_0	β_1	β_2
V	0.39 ± 0.03	-1.03 ± 0.25	3.50 ± 1.55
A_0	0.39 ± 0.05	-0.78 ± 0.26	2.41 ± 1.48
A_1	0.30 ± 0.03	0.48 ± 0.19	0.29 ± 0.65
A_3	0.25 ± 0.03	0.76 ± 0.20	0.71 ± 0.96
T_1	0.31 ± 0.03	-0.87 ± 0.19	2.75 ± 1.19
T_2	0.31 ± 0.03	0.58 ± 0.19	0.89 ± 0.71
T_3	0.68 ± 0.07	2.11 ± 0.46	4.94 ± 2.25

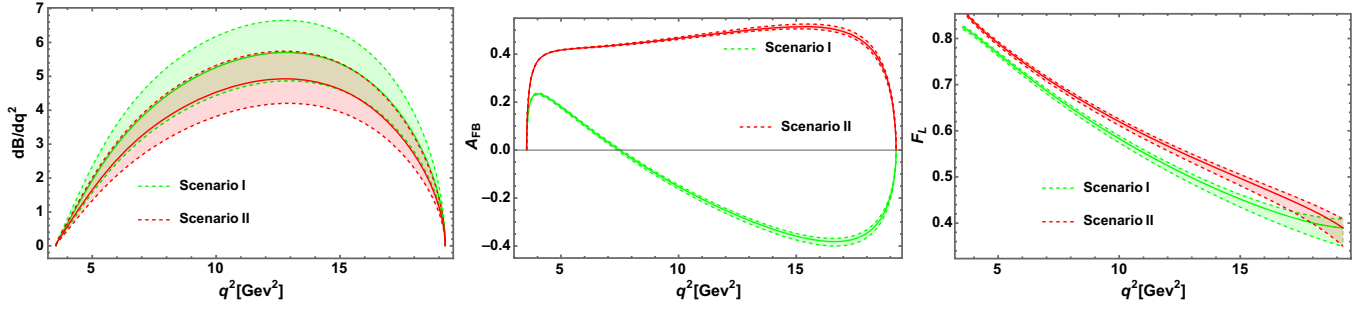


FIG. 2. Branching ratio (in units of 10^{-5}) (left), forward-backward asymmetry (middle), and lepton polarization (right) for $B \rightarrow K^* \tau^+ \mu^-$.

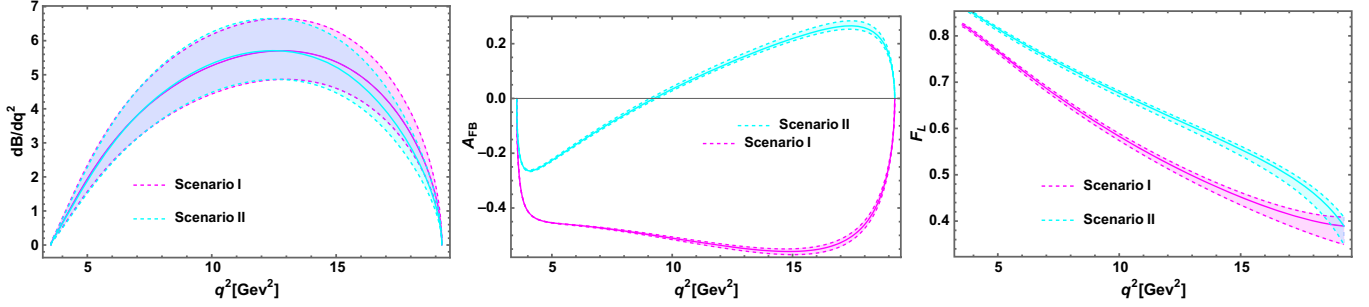


FIG. 3. Branching ratio (in units of 10^{-5}) (left), forward-backward asymmetry (middle), and lepton polarization (right) for $B \rightarrow K^* \tau^- \mu^+$.

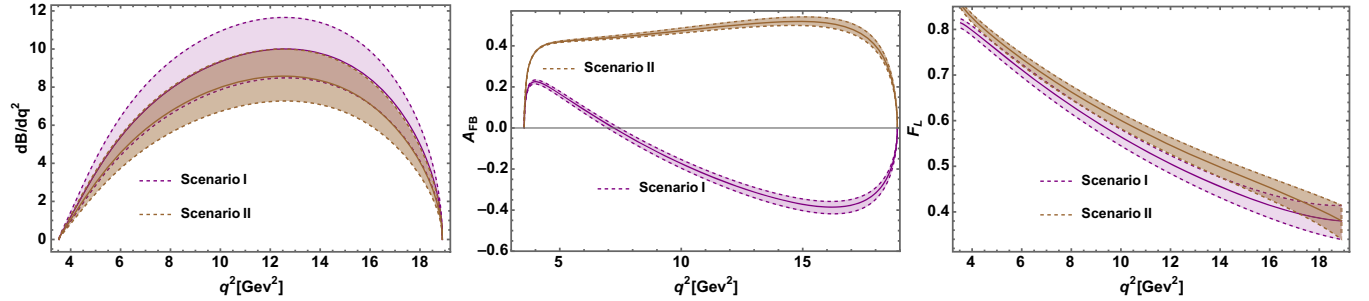


FIG. 4. Branching ratio (in units of 10^{-5}) (left), forward-backward asymmetry (middle), and lepton polarization (right) for $B \rightarrow \phi \tau^+ \mu^-$.

Scenario II (brown), each shown with their corresponding 1σ uncertainty bands. Similarly, Fig. 5 shows the corresponding observables for the $B \rightarrow \phi \tau^- \mu^+$ decay, with Scenario I illustrated in orange and Scenario II in blue, again including 1σ error bands.

Here, in the branching ratio plots, we can see that each mode of the decay process has almost similar q^2 dependencies irrespective of the charge of the heavier lepton, but in case of the $\tau^+ \mu^-$ channel the central value of Scenario I is higher than that of Scenario II. A_{FB} is particularly sensitive to the chiral structure of the underlying NP interactions. A consistent zero-crossing point near $q^2 \approx 9 \text{ GeV}^2$ appears in both plots, characteristic of the interplay between vector and axial-vector contributions in the decay dynamics. Scenario I demonstrates a significant positive asymmetry

for the $\tau^+ \mu^-$ channel at higher q^2 values (peaking near 0.4), while Scenario II of $\tau^- \mu^+$ remains predominantly negative. This divergence suggests potential new physics contributions. $\mathcal{F}_L$ in $B \rightarrow K^* \ell^+ \ell^-$ decays exhibits consistent behavior across Scenarios I and II, with values spanning $0.4 \leq F_L \leq 0.8$. Both scenarios show a monotonic decrease in F_L with increasing q^2 , suggesting a growing contribution from transverse amplitudes at higher dilepton invariant masses. The negligible difference between μ and τ final states indicates that F_L is less sensitive to lepton mass effects compared to observables like A_{FB} . Marginal deviations between scenarios at high q^2 ($\sim 15 \text{ GeV}^2$) may reflect theoretical uncertainties or new physics contributions to the helicity structure. Finally, Table IV summarizes the numerical values of the $\mathcal{B}$, A_{FB} , and $\mathcal{F}_L$ for the $B \rightarrow K^* \tau^\pm \mu^\mp$ LFV

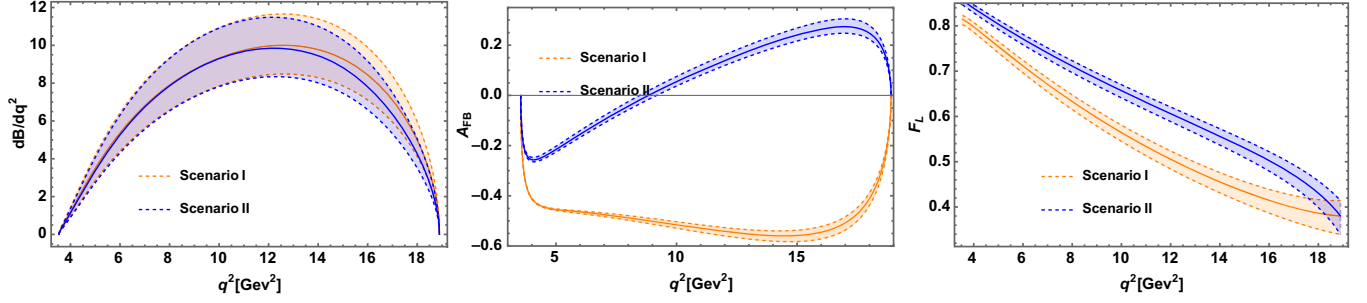


FIG. 5. Branching ratio (in units of 10^{-5}) (left), forward-backward asymmetry (middle), and lepton polarization (right) for $B \rightarrow \phi\tau^-\mu^+$.

decays. These results incorporate the upper limit constraints at the 90% confidence level for both scenarios.

IV. DECAY DISTRIBUTION OF $B \rightarrow K_2^*\ell_1\ell_2$

The NP Wilson coefficients $C_{V,A,S,P}^{(\ell)}$ encompass the short-distance physical aspects, while the long-distance

characteristics are integrated into the hadronic matrix elements of $B \rightarrow K_2^*$. In the context of the $B \rightarrow K_2^*$ transition, the hadronic matrix elements associated with the V and A currents can be described using four form factors: $V(q^2)$ and $A_{0,1,2}(q^2)$. These can be written as [97]

$$\begin{aligned} \langle K_2^*(k, \epsilon^*) | \bar{s}\gamma^\mu \gamma_5 b | \bar{B}(p) \rangle &= i(m_B + m_{K_2^*})A_1(q^2) \left[\epsilon_T^{*\mu} - \frac{\epsilon_T^* \cdot q}{q^2} q^\mu \right] + 2im_{K_2^*}A_0(q^2) \frac{\epsilon_T^* \cdot q}{q^2} q^\mu \\ &\quad - iA_2(q^2) \frac{\epsilon_T^* \cdot q}{m_B + m_{K_2^*}} \left[(p+k)^\mu - \frac{m_B^2 - m_{K_2^*}^2}{q^2} q^\mu \right], \\ \langle K_2^*(k, \epsilon^*) | \bar{s}\gamma^\mu b | \bar{B}(p) \rangle &= -\frac{2V(q^2)}{m_B + m_{K_2^*}} \epsilon^{\mu\nu\rho\sigma} \epsilon_{T\nu}^* p_\rho k_\sigma. \end{aligned} \quad (32)$$

However, by applying the equation of motion, it can be demonstrated that the matrix element for the $B \rightarrow K_2^*$ transition involving the scalar interaction ($\bar{s}b$) becomes null, while the one associated with the pseudoscalar interaction can be expressed as

$$\langle K_2^*(k, \epsilon^*) | \bar{s}\gamma_5 b | \bar{B}(p) \rangle = -\frac{2im_{K_2^*}A_0(q^2)}{m_b + m_s} (\epsilon_T^* \cdot q). \quad (33)$$

Here, p and k represent the four-momentum of the B and K_2^* mesons, respectively. The details regarding the polarization vector of K_2^* can be found in Appendix B. All the form factors have been computed using both the perturbative QCD approach [97] and the light cone QCD sum rule technique [98,99]. For this study, we have adopted the most up-to-date values of the form factors, as determined through the light cone QCD sum rule method, as presented in Ref. [98]. Within this method, each form factor can be expressed as follows:

$$F^{B \rightarrow K_2^*}(q^2) = \frac{1}{1 - q^2/m_{R,F}^2} \sum_{n=0}^1 \beta_n^F [z(q^2) - z(0)]^n, \quad (34)$$

where $z(q^2) = \frac{\sqrt{t_+ - s} - \sqrt{t_+ - t_0}}{\sqrt{t_+ - s} + \sqrt{t_+ - t_0}}$, $t_\pm = (m_B \pm m_{K_2^*})^2$, $t_0 = t_+ (1 - \sqrt{1 - t_-/t_+})$, and $m_{R,F}$ is the resonance mass associated with the quantum numbers of the corresponding form factor. The resonance masses employed in our

TABLE IV. Upper limits for the branching ratio, forward-backward asymmetry ($\mathcal{A}_{FB}$), and lepton polarization (F_L) of $B \rightarrow (K^*, \phi)\tau^\pm\mu^\mp$ LFV decays considering bounds are at the 90% CL.

Decay modes	Observables	Scenario I	Scenario II
$B \rightarrow K^*\tau^+\mu^-$	$\mathcal{B}$	$< 6.547 \times 10^{-5}$	$< 5.606 \times 10^{-5}$
	$\mathcal{A}_{FB}$	< -0.156	< 0.045
	$\mathcal{F}_L$	< 0.563	< 0.596
$B \rightarrow K^*\tau^-\mu^+$	$\mathcal{B}$	$< 6.514 \times 10^{-5}$	$< 6.415 \times 10^{-5}$
	$\mathcal{A}_{FB}$	< -0.490	< 0.055
	$\mathcal{F}_L$	< 0.561	< 0.642
$B \rightarrow \phi\tau^+\mu^-$	$\mathcal{B}$	$< 1.127 \times 10^{-5}$	$< 1.577 \times 10^{-5}$
	$\mathcal{A}_{FB}$	< -0.164	< 0.455
	$\mathcal{F}_L$	< 0.550	< 0.583
$B \rightarrow \phi\tau^-\mu^+$	$\mathcal{B}$	$< 1.123 \times 10^{-5}$	$< 1.087 \times 10^{-5}$
	$\mathcal{A}_{FB}$	< -0.491	< 0.066
	$\mathcal{F}_L$	< 0.557	< 0.627

TABLE V. Fits for $B \rightarrow K_2^*$ form factors using the LCSR approach.

Form factor	β_1	β_0
$V^{B \rightarrow K_2^*}$	$-0.90^{+0.37}_{-0.50}$	$0.22^{+0.11}_{-0.08}$
$A_0^{B \rightarrow K_2^*}$	$-1.23^{+0.23}_{-0.23}$	$0.30^{+0.01}_{-0.05}$
$A_1^{B \rightarrow K_2^*}$	$-0.46^{+0.19}_{-0.25}$	$0.19^{+0.09}_{-0.07}$
$A_2^{B \rightarrow K_2^*}$	$-0.40^{+0.23}_{-0.16}$	$0.11^{+0.05}_{-0.06}$

numerical calculations are given as

$$m_{R,V} = 5.412 \text{ GeV}, \quad m_{R,A_0} = 5.336 \text{ GeV},$$

$$m_{R,(A_1,A_3)} = 5.829 \text{ GeV}. \quad (35)$$

The fit parameters β_n^F are given in Table V. The three-body $B \rightarrow K_2^* \ell_1 \ell_2$ decay can be characterized using the leptonic

polar angle θ_ℓ and the squared leptonic mass, denoted as $q^2 = (p - k)^2$. We define the θ_ℓ angle as the angle formed by the ℓ_1 lepton concerning the rest frame of the dilepton system. In terms of these two variables, the two-fold differential decay distribution can be expressed as follows:

$$\frac{d^2\Gamma}{dq^2 d\cos\theta_\ell} = A(q^2) + B(q^2) \cos\theta_\ell + C(q^2) \cos^2\theta_\ell, \quad (36)$$

where

$$C(q^2) = \frac{3}{8} \beta_-^2 \beta_+^2 \{(|A_L^\parallel|^2 + |A_L^\perp|^2 - 2|A_L^0|^2) + (L \rightarrow R)\}, \quad (37)$$

$$B(q^2) = \frac{3}{2} \beta_+ \beta_- \left\{ \text{Re}[A_L^{\perp*} A_L^\parallel - (L \rightarrow R)] + \frac{m_+ m_-}{q^2} \text{Re}[A_L^0 A_L^t + (L \rightarrow R)] \right.$$

$$\left. + \frac{m_+}{\sqrt{q^2}} \text{Re}[A_S^*(A_L^0 + A_R^0)] - \frac{m_-}{\sqrt{q^2}} \text{Re}[A_{SP}^*(A_L^0 - A_R^0)] \right\}, \quad (38)$$

$$A(q^2) = \frac{3}{4} \left\{ \frac{1}{4} \left[\left(1 + \frac{m_-^2}{q^2}\right) \beta_+^2 + \left(1 + \frac{m_+^2}{q^2}\right) \beta_-^2 \right] (|A_L^\parallel|^2 + |A_L^\perp|^2 + (L \rightarrow R)) + \frac{1}{2} (\beta_+^2 + \beta_-^2) (|A_L^0|^2 + |A_R^0|^2) \right.$$

$$+ \frac{4m_1 m_2}{q^2} \text{Re}[A_L^0 A_R^0 + A_L^{\parallel*} A_R^\parallel + A_R^\perp A_L^{\perp*} - A_L^t A_R^{t*}]$$

$$+ \frac{1}{2} (\beta_+^2 + \beta_-^2 - 2\beta_-^2 \beta_+^2) (|A_L^t|^2 + |A_R^t|^2) + \frac{1}{2} (|A_{SP}|^2 \beta_-^2 + |A_S|^2 \beta_+^2)$$

$$\left. + \frac{2m_-}{\sqrt{q^2}} \beta_+^2 \text{Re}[A_S(A_L^t + A_R^t)^*] - \frac{2m_+}{\sqrt{q^2}} \beta_-^2 \text{Re}[A_{SP}(A_L^t - A_R^t)^*] \right\}. \quad (39)$$

Here, $m_\pm = (m_1 \pm m_2)$, $\beta_\pm = \sqrt{1 - \frac{(m_{\ell_1} \pm m_{\ell_2})^2}{q^2}}$, and A 's are the transversity amplitudes. Integrating Eq. (36) over θ_ℓ , we get the differential decay rate

$$\frac{d\Gamma}{dq^2} = 2 \left(A + \frac{C}{3} \right), \quad (40)$$

whereas the lepton forward-backward asymmetry is found to be

$$A_{\text{FB}}(q^2) = \frac{1}{d\Gamma/dq^2} \left(\int_0^1 d\cos\theta_\ell \frac{d^2\Gamma}{d\cos\theta_\ell dq^2} - \int_{-1}^0 d\cos\theta_\ell \frac{d^2\Gamma}{d\cos\theta_\ell dq^2} \right) = \frac{B}{2(A + C/3)}. \quad (41)$$

The $A_{\text{FB}}(q^2)$ in LFV processes is highly sensitive to the choice of hadronic form factor parametrization, as the observable is entirely governed by new physics operators. These contributions depend critically on the interference between vector, axial-vector, and tensor amplitudes, each constructed from distinct combinations of nonperturbative

form factors. Consequently, both the shape and potential zero-crossing of $A_{\text{FB}}(q^2)$ are influenced by the detailed q^2 dependence of these inputs. In our analysis, we employ LCSR form factors with the Ball-Zwicky dipole parametrization, which is valid for $q^2 \leq 12\text{--}14 \text{ GeV}^2$. However, alternative approaches, such as the Bourrely-Capri-

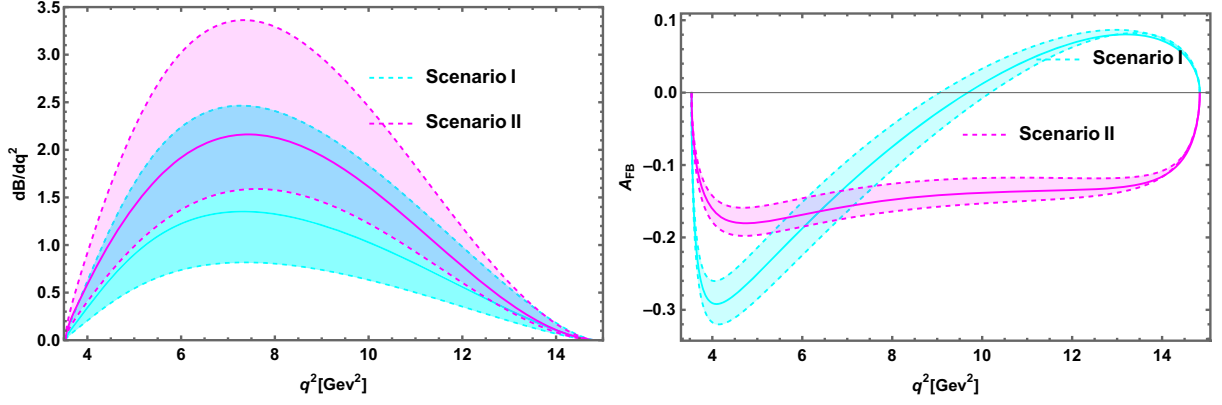


FIG. 6. Branching ratio (in units of 10^{-6}) (left) and forward-backward asymmetry (right) for $B \rightarrow K_2^* \tau^+ \mu^-$.

Lellouch z -expansion or lattice QCD inputs at high q^2 , can lead to quantitatively different predictions. These may affect not only the overall magnitude but also the shape and location of the zero asymmetry, if present. A complete treatment would require propagating form factor uncertainties, including correlations and parametrization dependence, which we defer to future dedicated studies.

A. Numerical analysis for $B \rightarrow K_2^* \ell_1 \ell_2$

In Fig. 6, we present the differential distributions for the $B \rightarrow K_2^* \tau^+ \mu^-$ decay. The plots display the $\mathcal{B}$ (left) and $\mathcal{A}_{FB}$ (right), highlighting the differences between Scenario I (cyan blue) and Scenario II (magenta), each shown with their corresponding 1σ uncertainty bands. Similarly, Fig. 7 shows the corresponding observables for the $B \rightarrow K_2^* \tau^- \mu^+$ decay, with Scenario I illustrated in purple and Scenario II in orange, again including 1σ error bands.

The branching ratios for the decays $B \rightarrow K_2^* \ell^+ \ell^-$ exhibit qualitatively similar behavior for both $\tau^\pm \mu^\mp$ channels. In both scenarios, a pronounced peak appears in the low- q^2 region, followed by a steep decrease. Scenario II yields a higher central value compared to Scenario I. The differential branching ratio plots indicate that the overall magnitude remains consistent between both decay modes. Regarding

the forward-backward asymmetry $\mathcal{A}_{FB}$, Scenario II predicts negative values throughout the full q^2 spectrum for both processes. In contrast, Scenario I shows a zero-crossing at $q^2 = 10 \text{ GeV}^2$ for the $B \rightarrow K_2^* \tau^+ \mu^-$ channel. Table VI summarizes the numerical results for both the branching ratios and forward-backward asymmetries in the lepton flavor-violating decays $B \rightarrow K_2^* \tau^+ \mu^-$ and $B \rightarrow K_2^* \tau^- \mu^+$, evaluated under the assumption of 90% confidence level upper limit constraints.

V. DECAY DISTRIBUTION OF $\Lambda_b \rightarrow \Lambda \ell_1^+ \ell_2^-$

To establish the kinematics of the decay process, we assume that the Λ_b particle is stationary, while the Λ and the dilepton pair move in opposite directions along the positive and negative z -axes, respectively. We denote the momenta of these particles as follows: p for the Λ_b , k for the Λ , q_1 for ℓ_1 , and q_2 for ℓ_2 . Additionally, we define s_p and s_k as the spins of the Λ_b and Λ , respectively, with respect to the z -axis in their respective rest frames. We also introduce two key kinematic variables: $q^\mu = q_1^\mu + q_2^\mu$ represents the four-momentum of the dilepton pair, and θ_ℓ represents the angle at which the lepton ℓ_1 is oriented relative to the z -axis in the rest frame of the dilepton system. The decay amplitudes can be written as

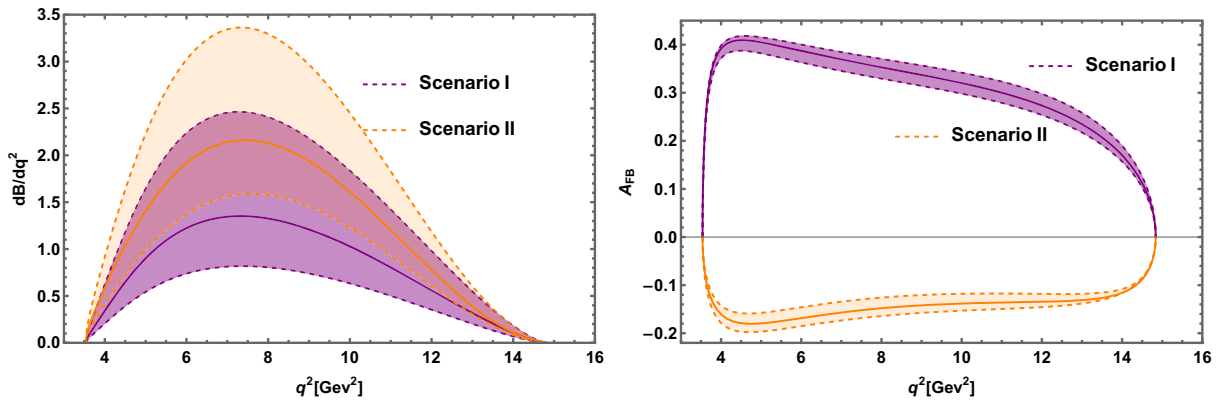


FIG. 7. Branching ratio (in units of 10^{-6}) (left) and forward-backward asymmetry (right) for $B \rightarrow K_2^* \tau^- \mu^+$.

TABLE VI. Upper limits for the branching ratio, forward-backward asymmetry ($\mathcal{A}_{FB}$), and lepton polarization ($\mathcal{F}_L$) of $B \rightarrow K_2^* \tau^\pm \mu^\mp$ LFV decays considering bounds are at the 90% CL.

Decay modes	Observables	Scenario I	Scenario II
$B \rightarrow K_2^* \tau^+ \mu^-$	$\mathcal{B}$	$< 9.131 \times 10^{-6}$	$< 1.416 \times 10^{-6}$
	$\mathcal{A}_{FB}$	< -0.056	< -0.142
$B \rightarrow K_2^* \tau^- \mu^+$	$\mathcal{B}$	$< 9.131 \times 10^{-6}$	$< 1.416 \times 10^{-6}$
	$\mathcal{A}_{FB}$	< 0.311	< -0.142

$$\mathcal{M}^{\lambda_2, \lambda_1}(s_p, s_k) = -\frac{V_{tb} V_{ts}^* \alpha_e}{2v^2} \sum_{i=L,R} \left[\sum_{\lambda} \eta_{\lambda} H_{VA,\lambda}^{i,s_p,s_k} L_{i,\lambda}^{\lambda_2, \lambda_1} + H_{SP}^{i,s_p,s_k} L_i^{\lambda_2, \lambda_1} \right]. \quad (42)$$

In this context, we have two sets of amplitudes: $H_{VA,\lambda}^{i,s_p,s_k}$ and H_{SP}^{i,s_p,s_k} . These amplitudes correspond to the vector and axial-vector (VA) operators as well as the scalar and pseudoscalar (SP) operators. Additionally, we have the leptonic helicity amplitudes denoted as $L_{i,\lambda}^{\lambda_2, \lambda_1}$ and $L_i^{\lambda_2, \lambda_1}$. Here, the subscript i represents the chiralities of the lepton current, and λ can take on values of t , $+1$, -1 , or 0 , which

correspond to the helicity states of the virtual gauge boson that decays into the dilepton pair. The subscripts $\lambda_{1,2}$ refer to the helicities of the leptons, and η_t is equal to 1, while $\eta_{\pm 1,0}$ is equal to -1 . For the detailed definitions and expressions of $H_{VA,\lambda}^{i,s_p,s_k}$ and H_{SP}^{i,s_p,s_k} in terms of Wilson coefficients and form factors, you can refer to [100]. In the existing literature, transversity amplitudes, denoted as $A_{\perp(\parallel)_1}^i$, $A_{\perp(\parallel)_0}^i$, $A_{S\perp(\parallel)}$, and $A_{P\perp(\parallel)}$, are frequently used instead of the hadronic helicity amplitudes. You can find the expressions for these transversity amplitudes in Appendix C according to the conventions outlined in [101]. The $L_i^{\lambda_2, \lambda_1}$ and $L_{i,\lambda}^{\lambda_2, \lambda_1}$ amplitudes are defined as

$$L_{L(R)}^{\lambda_2, \lambda_1} = \langle \bar{\ell}_2(\lambda_2) \ell_1(\lambda_1) | \bar{\ell}_2(1 \mp \gamma_5) \ell_1 | 0 \rangle, \\ L_{L(R),\lambda}^{\lambda_2, \lambda_1} = \bar{\epsilon}^\mu(\lambda) \langle \bar{\ell}_2(\lambda_2) \ell_1(\lambda_1) | \bar{\ell}_2 \gamma_\mu (1 \mp \gamma_5) \ell_1 | 0 \rangle. \quad (43)$$

Here, ϵ^μ represents the polarization vector of the virtual gauge boson that undergoes decay into the dilepton pair. The specific computations for the values of $L_{i,\lambda}^{\lambda_2, \lambda_1}$ and $L_i^{\lambda_2, \lambda_1}$ are provided in Appendix D. Utilizing these calculations, we determine the differential branching ratio for the decay process $\Lambda_b \rightarrow \Lambda \ell_1 \ell_2$ as

$$\frac{d\mathcal{B}}{dq^2 d\cos\theta_\ell} = \frac{3}{2} (K_{1ss}(1 - \cos^2\theta_\ell) + K_{1cc}\cos^2\theta_\ell + K_{1c}\cos\theta_\ell). \quad (44)$$

Each of the angular coefficients $K_{1ss,1cc,1c}$ can be written in the following way:

$$K_{1ss,1cc} = K_{1ss,1cc}^{VA} + K_{1ss,1cc}^{SP} + K_{1ss,1cc}^{int}. \quad (45)$$

The contributions from VA and SP operators are denoted as $K_{1ss,1cc,1c}^{VA}$ and $K_{1ss,1cc,1c}^{SP}$, and their interference terms are included in $K_{1ss,1cc,1c}^{int}$. In terms of the transversity amplitudes, the expressions for $K_{1ss,1cc,1c}^{VA}$ and $K_{1ss,1cc,1c}^{SP}$ are as follows:

$$K_{1ss}^{VA} = \frac{1}{4} (2|A_{\parallel 0}^R|^2 + |A_{\parallel 1}^R|^2 + 2|A_{\perp 0}^R|^2 + |A_{\perp 1}^R|^2 + \{R \leftrightarrow L\}) - \frac{m_+^2 + m_-^2}{4q^2} \\ \times [(|A_{\parallel 0}^R|^2 + |A_{\perp 0}^R|^2 + \{R \leftrightarrow L\}) - (|A_{\perp t}|^2 + \{\perp \leftrightarrow \parallel\})] + \frac{m_+^2 - m_-^2}{4q^2} \\ \times [2\text{Re}(A_{\perp 0}^R A_{\perp 0}^{*L} + A_{\perp 1}^R A_{\perp 1}^{*L} + \{\perp \leftrightarrow \parallel\})] - \frac{m_+^2 m_-^2}{4q^4} [(|A_{\parallel 1}^R|^2 + |A_{\perp 1}^R|^2 + \{R \leftrightarrow L\}) + 2|A_{\parallel t}|^2 + 2|A_{\perp t}|^2], \quad (46)$$

$$K_{1cc}^{VA} = \frac{1}{2} (|A_{\parallel 1}^R|^2 + |A_{\perp 1}^R|^2 + \{R \leftrightarrow L\}) + \frac{m_+^2 + m_-^2}{4q^2} [(|A_{\parallel 0}^R|^2 - |A_{\parallel 1}^R|^2 + |A_{\perp 0}^R|^2 \\ - |A_{\perp 1}^R|^2 + \{R \leftrightarrow L\}) + (|A_{\perp t}|^2 + |A_{\parallel t}|^2)] + \frac{m_+^2 - m_-^2}{4q^2} [2\text{Re}(A_{\perp 0}^R A_{\perp 0}^{*L} + A_{\perp 1}^R A_{\perp 1}^{*L} + \{\perp \leftrightarrow \parallel\})] \\ - \frac{m_+^2 m_-^2}{2q^4} [(|A_{\parallel 0}^R|^2 + |A_{\perp 0}^R|^2 + \{R \leftrightarrow L\}) + |A_{\parallel t}|^2 + |A_{\perp t}|^2], \quad (47)$$

$$K_{1c}^{\text{VA}} = -\beta_\ell\beta'_\ell(A_{\perp 1}^R A_{\parallel 1}^{*R} - \{R \leftrightarrow L\}) + \beta_\ell\beta'_\ell \frac{m_+ + m_-}{q^2} \text{Re}(A_{\parallel 0}^L A_{\parallel t}^{*} + A_{\perp 0}^L A_{\perp t}^{*}), \quad (48)$$

$$K_{1ss}^{\text{SP}} = \frac{1}{4}(|A_{S\perp}|^2 + |A_{P\perp}|^2 + \{\perp \leftrightarrow \parallel\}) - \frac{m_+^2}{4q^2}(|A_{S\parallel}|^2 + |A_{S\perp}|^2) - \frac{m_-^2}{4q^2}(|A_{P\parallel}|^2 + |A_{P\perp}|^2), \quad (49)$$

$$K_{1cc}^{\text{SP}} = \frac{1}{4}(|A_{P\perp}|^2 + |A_{S\perp}|^2 + \{\perp \leftrightarrow \parallel\}) - \frac{m_+^2}{4q^2}(|A_{S\parallel}|^2 + |A_{S\perp}|^2) - \frac{m_-^2}{4q^2}(|A_{P\parallel}|^2 + |A_{P\perp}|^2), \quad (50)$$

$$K_{1c}^{\text{SP}} = 0, \quad (51)$$

$$K_{1ss}^{\text{int}} = \frac{m_+}{2\sqrt{q^2}} \text{Re}(A_{\parallel t} A_{P\parallel}^{*} + A_{\perp t} A_{P\perp}^{*}) + \frac{m_-}{2\sqrt{q^2}} \text{Re}(A_{\parallel t} A_{S\parallel}^{*} + A_{\perp t} A_{S\perp}^{*}) \\ - \frac{m_+^2 m_-}{2q^2 \sqrt{q^2}} \text{Re}(A_{\parallel t} A_{S\parallel}^{*} + A_{\perp t} A_{S\perp}^{*}) - \frac{m_+ m_-^2}{2q^2 \sqrt{q^2}} \text{Re}(A_{\parallel t} A_{P\parallel}^{*} + A_{\perp t} A_{P\perp}^{*}), \quad (52)$$

$$K_{1cc}^{\text{int}} = \frac{m_+}{2\sqrt{q^2}} \text{Re}(A_{\parallel t} A_{P\parallel}^{*} + A_{\perp t} A_{P\perp}^{*}) + \frac{m_-}{2\sqrt{q^2}} \text{Re}(A_{\parallel t} A_{S\parallel}^{*} + A_{\perp t} A_{S\perp}^{*}) \\ - \frac{m_+^2 m_-}{2q^2 \sqrt{q^2}} \text{Re}(A_{\parallel t} A_{S\parallel}^{*} + A_{\perp t} A_{S\perp}^{*}) - \frac{m_+ m_-^2}{2q^2 \sqrt{q^2}} \text{Re}(A_{\parallel t} A_{P\parallel}^{*} + A_{\perp t} A_{P\perp}^{*}), \quad (53)$$

$$K_{1c}^{\text{int}} = \frac{\beta_\ell\beta'_\ell}{2\sqrt{q^2}} \text{Re}(A_{S\parallel} A_{\parallel 0}^{L*} + A_{S\perp} A_{\perp 0}^{L*} + A_{S\parallel} A_{\parallel 0}^{R*} + A_{S\perp} A_{\perp 0}^{R*}) \\ + \frac{\beta_\ell\beta'_\ell}{2\sqrt{q^2}} \text{Re}(A_{P\parallel} A_{\parallel 0}^{L*} + A_{P\perp} A_{\perp 0}^{L*} - A_{P\parallel} A_{\parallel 0}^{R*} - A_{P\perp} A_{\perp 0}^{R*}). \quad (54)$$

We have introduced the notations $m_\pm = m_1 \pm m_2$, where m_1 and m_2 represent the masses of ℓ_1 and ℓ_2 , respectively. The relative importance of the terms in Eqs. (51) and (52) depends sensitively on both the lepton masses and the chosen q^2 range. When one of the leptons is a τ , the terms proportional to m_τ (second line) typically dominate due to the mass enhancement, especially at higher q^2 . Conversely, for light leptons (e, μ), the first line dominates at low q^2 due to kinematic suppression of the heavy lepton terms. This behavior does, in principle, allow one to disentangle contributions based on the q^2 region and lepton flavor structure, particularly in differential distributions. This feature may be useful in guiding future experimental analyses targeting LFV transitions with asymmetric lepton masses. The coefficients $\beta_\ell^{(\prime)}$ are specified in Appendix C. Based on the differential decay distribution, we establish two observables, as outlined in [101]: the differential branching ratio is

$$\frac{d\mathcal{B}}{dq^2} = 2K_{1ss} + K_{1cc} \quad (55)$$

and the forward-backward asymmetry is

$$A_{\text{FB}}^\ell = \frac{3}{2} \frac{K_{1c}}{K_{1ss} + K_{1cc}}. \quad (56)$$

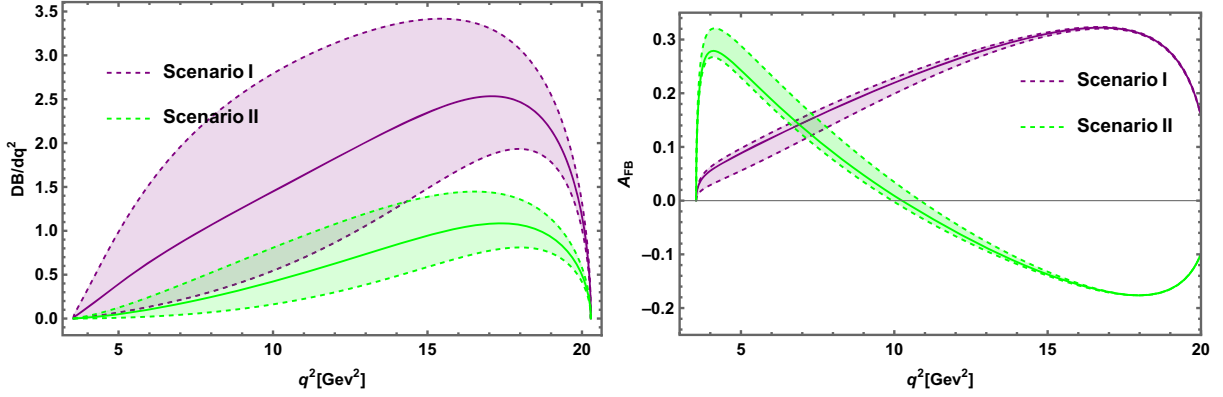
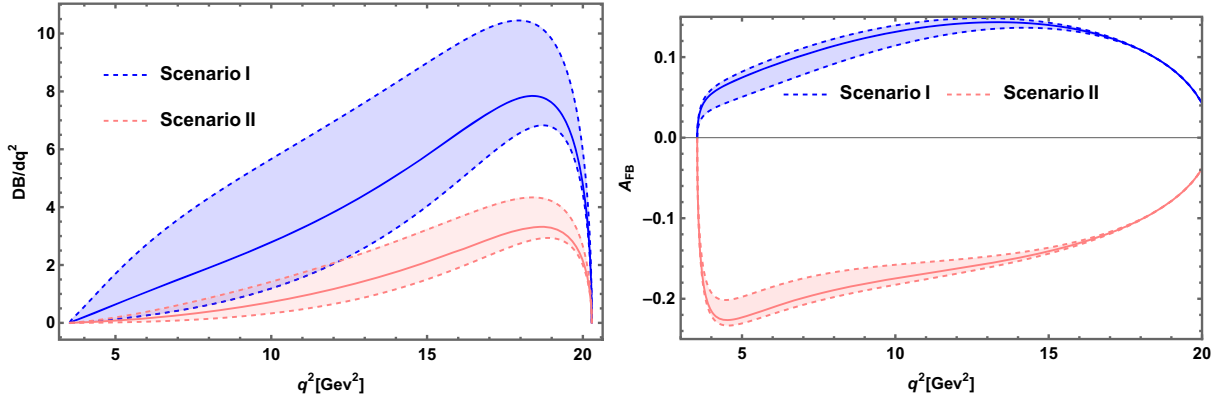
The available phase space in the dilepton invariant mass squared q^2 is

$$(m_1 + m_2)^2 \leq q^2 \leq (m_{\Lambda_b} - m_\Lambda)^2. \quad (57)$$

A. Numerical analysis for $\Lambda_b \rightarrow \Lambda \ell_1 \ell_2$

Given that the mesonic decays $B^+ \rightarrow K^+ \ell_1 \ell_2$ and $\bar{B}_s \rightarrow \ell_1 \ell_2$ arise from the same underlying quark-level transition as the baryonic mode $\Lambda_b \rightarrow \Lambda \ell_1 \ell_2$, we utilize the constrained NP parameter space, obtained from these mesonic channels and illustrated in Fig. 1, to study the corresponding baryonic decays. Specifically, we investigate the q^2 dependence of the differential branching ratio and the forward-backward asymmetry for the LFV processes $\Lambda_b \rightarrow \Lambda \tau^\pm \mu^\mp$.

Figure 8 displays the q^2 distributions for the $\Lambda_b \rightarrow \Lambda \tau^+ \mu^-$ decay. The left panel presents the differential branching ratio $d\mathcal{B}/dq^2$, while the right shows the forward-backward asymmetry $\mathcal{A}_{\text{FB}}(q^2)$, comparing Scenario I (cyan blue) and Scenario II (magenta), each with their associated 1σ uncertainty bands. Similarly, Fig. 9 presents the

FIG. 8. Branching ratio (in units of 10^{-6}) (left) and forward-backward asymmetry (right) for $\Lambda_b \rightarrow \Lambda \tau^+ \mu^-$.FIG. 9. Branching ratio (in units of 10^{-6}) (left) and forward-backward asymmetry (right) for $\Lambda_b \rightarrow \Lambda \tau^- \mu^+$.

corresponding observables for the $\Lambda_b \rightarrow \Lambda \tau^- \mu^+$ decay, where Scenario I is shown in purple and Scenario II in orange.

The branching ratios of the LFV decays $\Lambda_b \rightarrow \Lambda \tau^\pm \mu^\mp$ demonstrate substantial deviations, reaching values up to $\mathcal{O}(10^{-5})$, thereby serving as clear signatures of LFV. In the $\tau^+ \mu^-$ mode, Scenario I dominates at low q^2 , whereas Scenario II becomes prominent at intermediate q^2 , suggesting contributions from heavier particles. A gradual increase in the branching ratio is observed with increasing q^2 , peaking around $q^2 \sim 3.5$ in Scenario II. The divergence between the two scenarios becomes more noticeable

beyond $q^2 \approx 10 \text{ GeV}^2$, indicating a strong dependence on the NP operator structure.

The $\mathcal{A}_{\text{FB}}(q^2)$ also reveals notable differences. For both decay modes, Scenario I exhibits a consistently positive asymmetry that peaks around $q^2 \sim 10\text{--}12 \text{ GeV}^2$ before tapering off at the kinematic extremes. In contrast, Scenario II shows a zero-crossing behavior in the $\tau^+ \mu^-$ decay and remains entirely negative in the $\tau^- \mu^+$ mode. These distinctions emphasize the asymmetry's sensitivity to the nature and chirality of the underlying NP interactions, with intermediate q^2 regions being particularly informative.

The quantitative results for the integrated branching ratios and forward-backward asymmetries are summarized in Table VII. These values incorporate the upper bound constraints on LFV branching ratios at the 90% confidence level and offer a detailed benchmark for future experimental investigations into baryonic LFV processes.

VI. SUMMARY

In this work, we have investigated a range of LFV decays: $B \rightarrow K^* \ell_1 \ell_2$, $B \rightarrow \phi \ell_1 \ell_2$, $B \rightarrow K_2^* \ell_1 \ell_2$, and $\Lambda_b \rightarrow \Lambda \ell_1 \ell_2$. Our study considers contributions from

TABLE VII. Upper limits for the branching ratio, forward-backward asymmetry ($\mathcal{A}_{\text{FB}}$), and lepton polarization ($\mathcal{F}_{\mathcal{L}}$) of $\Lambda_b \rightarrow \Lambda \tau^\pm \mu^\mp$ LFV decays considering bounds are at the 90% CL.

Decay modes	Observables	Scenario I	Scenario II
$\Lambda_b \rightarrow \Lambda \tau^+ \mu^-$	$\mathcal{B}$	$< 2.63 \times 10^{-5}$	$< 1.673 \times 10^{-5}$
	$\mathcal{A}_{\text{FB}}$	< 0.223	< -0.012
$\Lambda_b \rightarrow \Lambda \tau^- \mu^+$	$\mathcal{B}$	$< 6.583 \times 10^{-5}$	$< 9.631 \times 10^{-5}$
	$\mathcal{A}_{\text{FB}}$	< 0.111	< -0.160

vector, axial-vector, scalar, and pseudoscalar NP operators and examines two distinct scenarios under these assumptions. We focus on the double-differential decay distribution in terms of the dilepton invariant mass squared (q^2) and the lepton angle (θ_ℓ), from which we extract observables, such as the differential branching ratio, forward-backward asymmetry, and lepton polarization. The analysis is performed in a model-independent framework, allowing us to set upper bounds on the branching fractions and $\mathcal{A}_{FB}$ at the 90% confidence level, with potential sensitivity for experimental verification at LHCb. Our findings indicate that the branching ratios for the decays $B \rightarrow (K^*, \phi)\tau^\pm\mu^\mp$ are of the order $\mathcal{O}(10^{-5})$, with negligible difference between charge conjugate channels. For the $B \rightarrow K_2^*\tau^\pm\mu^\mp$ modes, the branching ratio is smaller, at $\mathcal{O}(10^{-6})$. We also identify the zero-crossing points of $\mathcal{A}_{FB}$ in the $B \rightarrow (K^*, \phi, K_2^*)\mu^+\tau^-$ modes. In the baryonic sector, the branching ratios for $\Lambda_b \rightarrow \Lambda\tau^+\mu^-$ and $\Lambda_b \rightarrow \Lambda\tau^-\mu^+$ are estimated to be 2.63×10^{-5} and 6.58×10^{-5} , respectively. Notably, the forward-backward asymmetry exhibits a zero-crossing only in the $\tau^+\mu^-$ mode, suggesting potential CP -violating or chiral NP effects. This charge asymmetry implies non-trivial structures in the LFV couplings, offering a promising window into distinguishing NP interactions by their lepton flavor and charge dependence. In summary, our analysis underscores the value of LFV observables as sensitive probes of new physics and highlights several decay channels as key targets for ongoing and future experimental investigations.

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DATA AVAILABILITY

The data are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX A: ANGULAR CONVENTIONS AND KINEMATICS

Our angular conventions for the decay $\bar{B} \rightarrow \bar{K}^*(\rightarrow K^-\pi^+)$ $\ell_1^-\ell_2^+$ are summarized in Fig. 10. In the B rest frame, the leptonic and hadronic four-vectors are defined by $q^\mu = (q_0, 0, 0, q_z)$ and $k^\mu = (k_0, 0, 0, -q_z)$, where

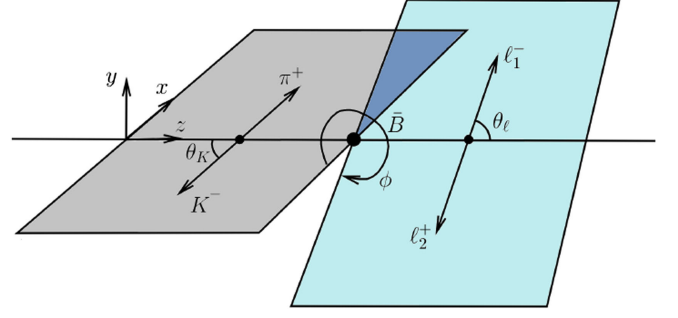


FIG. 10. Angular conventions for the decay $\bar{B} \rightarrow \bar{K}^*\ell_1^-\ell_2^+$.

$$q_0 = \frac{m_B^2 + q^2 - m_{K^*}^2}{2m_B}, \quad k_0 = \frac{m_B^2 + m_{K^*}^2 - q^2}{2m_B}, \quad \text{and} \quad q_z = \frac{\lambda^{1/2}(m_B, m_{K^*}, \sqrt{q^2})}{2m_B}. \quad (\text{A1})$$

In the dilepton rest frame, the leptonic four-vectors read

$$p_1^\mu = (E_\alpha, |p_\ell| \sin\theta_\ell \cos\phi, -|p_\ell| \sin\theta_\ell \sin\phi, |p_\ell| \cos\theta_\ell), \quad (\text{A2})$$

$$p_2^\mu = (E_\beta, -|p_\ell| \sin\theta_\ell \cos\phi, |p_\ell| \sin\theta_\ell \sin\phi, -|p_\ell| \cos\theta_\ell), \quad (\text{A3})$$

where

$$E_1 = \frac{q^2 + m_1^2 - m_2^2}{2\sqrt{q^2}}, \quad E_2 = \frac{q^2 - m_1^2 + m_2^2}{2\sqrt{q^2}}, \quad \text{and} \quad |p_\ell| = \frac{\lambda^{1/2}(q^2, m_1^2, m_2^2)}{2m_B}. \quad (\text{A4})$$

In the same way, one can write in the K^* rest frame

$$p_K^\mu = (E_K, -|p_K| \sin\theta_K, 0, -|p_K| \cos\theta_K), \quad (\text{A5})$$

$$p_\pi^\mu = (E_\pi, +|p_K| \sin\theta_K, 0, +|p_K| \cos\theta_K), \quad (\text{A6})$$

where E_K , E_π , and $|p_K|$ are given by similar expressions.

APPENDIX B: POLARIZATION OF K_2^*

The polarization $\epsilon^{\mu\nu}(n)$ of tensor meson K_2^* , which has four-momentum $(k_0, 0, 0, \vec{k})$, can be written in terms of the spin-1 polarization vectors [102]

$$\begin{aligned}
\epsilon_{\mu\nu}(\pm 2) &= \epsilon_\mu(\pm 1)\epsilon_\nu(\pm 1), \\
\epsilon_{\mu\nu}(\pm 1) &= \frac{1}{\sqrt{2}}[\epsilon_\nu(\pm)\epsilon_\mu(0) + \epsilon_\nu(\pm)\epsilon_\mu(0)], \\
\epsilon_{\mu\nu}(0) &= \frac{1}{\sqrt{6}}[\epsilon_\mu(+)\epsilon_\nu(-) + \epsilon_\nu(+)\epsilon_\mu(-)] + \sqrt{\frac{2}{3}}\epsilon_\mu(0)\epsilon_\nu(0),
\end{aligned} \tag{B1}$$

where the spin-1 polarization vectors are defined as

$$\epsilon_\mu(0) = \frac{1}{m_{K_2^*}}(\vec{k}_z, 0, 0, k_0), \quad \epsilon_\mu(\pm) = \frac{1}{\sqrt{2}}(0, 1, \pm i, 0). \tag{B2}$$

We study $B \rightarrow K_2^* \ell_1 \ell_2$ decay, where we have two leptons in the final state, so in this case, the $n = \pm 2$ helicity state of the K_2^* is not realized. Therefore, a new polarization vector is introduced [99],

$$\epsilon_{T\mu}(h) = \frac{\epsilon_{\mu\nu}P^\nu}{m_B}. \tag{B3}$$

The explicit expressions of polarization vectors are

$$\epsilon_{T\mu}(\pm 1) = \frac{1}{m_B} \frac{1}{\sqrt{2}} \epsilon(0) \cdot p \epsilon_\mu(\pm) = \frac{\sqrt{\lambda}}{\sqrt{8}m_B m_{K_2^*}} \epsilon_\mu(\pm), \tag{B4}$$

$$\epsilon_{T\mu}(0) = \frac{1}{m_B} \sqrt{\frac{2}{3}} \epsilon(0) \cdot p \epsilon_\mu(0) = \frac{\sqrt{\lambda}}{\sqrt{6}m_B m_{K_2^*}} \epsilon_\mu(0), \tag{B5}$$

where $\lambda(m_B^2, m_{K_2^*}^2, q^2) = m_B^4 + m_{K_2^*}^4 + q^4 - 2(m_B^2 m_{K_2^*}^2 + m_B^2 q^2 + m_{K_2^*}^2 q^2)$ is the usual Kallen function. On the other hand, the virtual gauge boson can have three types of polarization states, longitudinal, transverse, and timelike, which have the following components:

$$\begin{aligned}
\epsilon_V^\mu(0) &= \frac{1}{\sqrt{q^2}}(-|\vec{q}_z|, 0, 0, -q_0), \quad \epsilon_V^\mu(\pm) = \frac{1}{\sqrt{2}}(0, 1, \pm i, 0), \\
\epsilon_V^\mu(t) &= \frac{1}{\sqrt{q^2}}(q_0, 0, 0, q_z),
\end{aligned} \tag{B6}$$

where $q^\mu = (q_0, 0, 0, q_z)$ is the four-momentum of the gauge boson.

APPENDIX C: TRANSVERSITY AMPLITUDES

Corresponding to the effective Hamiltonian (3), the expressions of the transversity amplitudes read [101]

$$A_{\perp 1}^{L,(R)} = -\sqrt{2}N(f_{\perp}^V \sqrt{2s_-} \mathcal{C}_{VA+}^{L,(R)}), \tag{C1}$$

$$A_{\parallel 1}^{L,(R)} = \sqrt{2}N(f_{\perp}^A \sqrt{2s_+} \mathcal{C}_{VA-}^{L,(R)}), \tag{C2}$$

$$A_{\perp 0}^{L,(R)} = \sqrt{2}N\left(f_0^V(m_{\Lambda_b} + m_{\Lambda})\sqrt{\frac{s_-}{q^2}}\mathcal{C}_{VA+}^{L,(R)}\right), \tag{C3}$$

$$A_{\parallel 0}^{L,(R)} = -\sqrt{2}N\left(f_0^A(m_{\Lambda_b} - m_{\Lambda})\sqrt{\frac{s_+}{q^2}}\mathcal{C}_{VA-}^{L,(R)}\right), \tag{C4}$$

$$A_{\perp t} = -2\sqrt{2}Nf_t^V(m_{\Lambda_b} - m_{\Lambda})\sqrt{\frac{s_+}{q^2}}(\mathcal{C}_A + \mathcal{C}'_A), \tag{C5}$$

$$A_{\parallel t} = 2\sqrt{2}Nf_t^A(m_{\Lambda_b} + m_{\Lambda})\sqrt{\frac{s_-}{q^2}}(\mathcal{C}_A - \mathcal{C}'_A). \tag{C6}$$

Here, the normalization constant $N(q^2)$ is given by

$$\begin{aligned}
N(q^2) &= \frac{V_{tb}V_{ts}^*\alpha_e}{\sqrt{2}v^2} \sqrt{\tau_{\Lambda_b} \frac{q^2 \sqrt{\lambda(m_{\Lambda_b}^2, m_{\Lambda}^2, q^2)}}{2^{15}m_{\Lambda_b}^3 \pi^5} \beta_\ell \beta'_\ell}, \\
\beta_\ell &= \sqrt{1 - \frac{(m_1 + m_2)^2}{q^2}}, \quad \beta'_\ell = \sqrt{1 - \frac{(m_1 - m_2)^2}{q^2}},
\end{aligned} \tag{C7}$$

where $\lambda(a, b, c) = a^2 + b^2 + c^2 - 2(ab + bc + ca)$, and the Wilson coefficients are

$$\mathcal{C}_{VA,+}^{L,(R)} = (\mathcal{C}_V \mp \mathcal{C}_A) + (\mathcal{C}'_V \mp \mathcal{C}'_A), \tag{C8}$$

$$\mathcal{C}_{VA,-}^{L,(R)} = (\mathcal{C}_V \mp \mathcal{C}_A) - (\mathcal{C}'_V \mp \mathcal{C}'_A). \tag{C9}$$

The transversity amplitudes corresponding to the SP operators are [101]

$$A_{S\perp} = 2\sqrt{2}Nf_t^V \frac{m_{\Lambda_b} - m_{\Lambda}}{m_b} \sqrt{s_+}(\mathcal{C}_S + \mathcal{C}'_S), \tag{C10}$$

$$A_{S\parallel} = -2\sqrt{2}Nf_t^A \frac{m_{\Lambda_b} + m_{\Lambda}}{m_b} \sqrt{s_-}(\mathcal{C}_S - \mathcal{C}'_S), \tag{C11}$$

$$A_{P\perp} = -2\sqrt{2}Nf_t^V \frac{m_{\Lambda_b} - m_{\Lambda}}{m_b} \sqrt{s_+}(\mathcal{C}_P + \mathcal{C}'_P), \tag{C12}$$

$$A_{P\parallel} = 2\sqrt{2}Nf_t^A \frac{m_{\Lambda_b} + m_{\Lambda}}{m_b} \sqrt{s_-}(\mathcal{C}_P - \mathcal{C}'_P). \tag{C13}$$

APPENDIX D: SPINORS IN DILEPTON REST FRAME

We assume that the lepton ℓ_2^- is negatively charged and has four-momentum $q_2^\mu = (E_1, \vec{q})$, while ℓ_1^+ is positively charged and has four-momentum $q_1^\mu = (E_1, -\vec{q})$,

$$q_1^\mu|_{2\ell} = (E_2, -|q_{2\ell}|\sin\theta_\ell, 0, -|q_{2\ell}|\cos\theta_\ell), \quad (\text{D1})$$

$$q_2^\mu|_{2\ell} = (E_1, |q_{2\ell}|\sin\theta_\ell, 0, |q_{2\ell}|\cos\theta_\ell), \quad (\text{D2})$$

with

$$|q_{2\ell}| = \frac{\lambda^{1/2}(q^2, m_1^2, m_2^2)}{2\sqrt{q^2}}, \quad E_1 = \frac{q^2 + m_1^2 - m_2^2}{2\sqrt{q^2}},$$

$$E_2 = \frac{q^2 + m_2^2 - m_1^2}{2\sqrt{q^2}}. \quad (\text{D3})$$

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

$$\bar{u}_{\ell_2}(1 \mp \gamma_5)v_{\ell_1}, \quad \bar{e}^\mu(\lambda)\bar{u}_{\ell_2}\gamma_\mu(1 \mp \gamma_5)v_{\ell_1}. \quad (\text{D4})$$

With these choices of lepton spinors, we get the following expressions of the lepton helicity amplitudes $L_{L(R)}^{\lambda_2, \lambda_1}$ and $L_{L(R), \lambda}^{\lambda_2, \lambda_1}$:

$$L_L^{+\frac{1}{2}+\frac{1}{2}} = \sqrt{q^2}(\beta'_\ell + \beta_\ell), \quad L_L^{+\frac{1}{2}-\frac{1}{2}} = 0, \quad L_L^{-\frac{1}{2}+\frac{1}{2}} = 0, \quad L_L^{-\frac{1}{2}-\frac{1}{2}} = \sqrt{q^2}(\beta'_\ell - \beta_\ell), \quad (\text{D5})$$

$$L_R^{+\frac{1}{2}+\frac{1}{2}} = -\sqrt{q^2}(\beta'_\ell - \beta_\ell), \quad L_R^{+\frac{1}{2}-\frac{1}{2}} = 0, \quad L_R^{-\frac{1}{2}+\frac{1}{2}} = 0, \quad L_R^{-\frac{1}{2}-\frac{1}{2}} = -\sqrt{q^2}(\beta'_\ell + \beta_\ell), \quad (\text{D6})$$

$$L_{L,+1}^{+\frac{1}{2}+\frac{1}{2}} = \frac{1}{\sqrt{2}}[m_1(\beta'_\ell + \beta_\ell) + m_2(\beta'_\ell - \beta_\ell)], \quad L_{L,+1}^{+\frac{1}{2}-\frac{1}{2}} = -\sqrt{\frac{q^2}{2}}(\beta'_\ell - \beta_\ell)(1 - \cos\theta_\ell),$$

$$L_{L,+1}^{-\frac{1}{2}+\frac{1}{2}} = -\frac{1}{\sqrt{2}}[m_1(\beta'_\ell - \beta_\ell) + m_2(\beta'_\ell + \beta_\ell)], \quad L_{L,+1}^{-\frac{1}{2}-\frac{1}{2}} = \sqrt{\frac{q^2}{2}}(\beta'_\ell + \beta_\ell)(1 + \cos\theta_\ell), \quad (\text{D7})$$

$$L_{R,+1}^{+\frac{1}{2}+\frac{1}{2}} = \frac{1}{\sqrt{2}}[m_1(\beta'_\ell - \beta_\ell) + m_2(\beta'_\ell + \beta_\ell)], \quad L_{R,+1}^{+\frac{1}{2}-\frac{1}{2}} = -\sqrt{\frac{q^2}{2}}(\beta'_\ell + \beta_\ell)(1 - \cos\theta_\ell),$$

$$L_{R,+1}^{-\frac{1}{2}+\frac{1}{2}} = -\frac{1}{\sqrt{2}}[m_1(\beta'_\ell + \beta_\ell) + m_2(\beta'_\ell - \beta_\ell)], \quad L_{R,+1}^{-\frac{1}{2}-\frac{1}{2}} = \sqrt{\frac{q^2}{2}}(\beta'_\ell - \beta_\ell)(1 + \cos\theta_\ell), \quad (\text{D8})$$

$$L_{L,-1}^{+\frac{1}{2}+\frac{1}{2}} = -\frac{1}{\sqrt{2}}[m_1(\beta'_\ell + \beta_\ell) + m_2(\beta'_\ell - \beta_\ell)], \quad L_{L,-1}^{+\frac{1}{2}-\frac{1}{2}} = -\sqrt{\frac{q^2}{2}}(\beta'_\ell - \beta_\ell)(1 + \cos\theta_\ell),$$

$$L_{L,-1}^{-\frac{1}{2}+\frac{1}{2}} = \frac{1}{\sqrt{2}}[m_1(\beta'_\ell - \beta_\ell) + m_2(\beta'_\ell + \beta_\ell)], \quad L_{L,-1}^{-\frac{1}{2}-\frac{1}{2}} = \sqrt{\frac{q^2}{2}}(\beta'_\ell + \beta_\ell)(1 - \cos\theta_\ell), \quad (\text{D9})$$

$$L_{R,-1}^{+\frac{1}{2}+\frac{1}{2}} = -\frac{1}{\sqrt{2}}[m_1(\beta'_\ell - \beta_\ell) + m_2(\beta'_\ell + \beta_\ell)], \quad L_{R,-1}^{+\frac{1}{2}-\frac{1}{2}} = \sqrt{\frac{q^2}{2}}(\beta'_\ell + \beta_\ell)(1 + \cos\theta_\ell),$$

$$L_{R,-1}^{-\frac{1}{2}+\frac{1}{2}} = \frac{1}{\sqrt{2}}[m_1(\beta'_\ell + \beta_\ell) + m_2(\beta'_\ell - \beta_\ell)], \quad L_{R,-1}^{-\frac{1}{2}-\frac{1}{2}} = \sqrt{\frac{q^2}{2}}(\beta'_\ell - \beta_\ell)(1 - \cos\theta_\ell), \quad (\text{D10})$$

$$L_{L,0}^{+\frac{1}{2}+\frac{1}{2}} = -[m_1(\beta'_\ell + \beta_\ell) - m_2(\beta'_\ell - \beta_\ell)]\cos\theta_\ell, \quad L_{L,0}^{+\frac{1}{2}-\frac{1}{2}} = \sqrt{q^2}(\beta'_\ell - \beta_\ell)\cos\theta_\ell,$$

$$L_{L,0}^{-\frac{1}{2}+\frac{1}{2}} = [m_1(\beta'_\ell - \beta_\ell) + m_2(\beta'_\ell + \beta_\ell)]\cos\theta_\ell, \quad L_{L,0}^{-\frac{1}{2}-\frac{1}{2}} = \sqrt{q^2}(\beta'_\ell + \beta_\ell)\cos\theta_\ell, \quad (\text{D11})$$

$$L_{R,0}^{+\frac{1}{2}+\frac{1}{2}} = -[m_1(\beta'_\ell - \beta_\ell) + m_2(\beta'_\ell + \beta_\ell)]\cos\theta_\ell, \quad L_{R,0}^{+\frac{1}{2}-\frac{1}{2}} = \sqrt{q^2}(\beta'_\ell + \beta_\ell)\cos\theta_\ell,$$

$$L_{R,0}^{-\frac{1}{2}+\frac{1}{2}} = [m_1(\beta'_\ell + \beta_\ell) + m_2(\beta'_\ell - \beta_\ell)]\cos\theta_\ell, \quad L_{R,0}^{-\frac{1}{2}-\frac{1}{2}} = \sqrt{q^2}(\beta'_\ell - \beta_\ell), \quad (\text{D12})$$

$$\begin{aligned}
L_{L,0}^{+\frac{1}{2}+\frac{1}{2}} &= [m_1(\beta'_\ell + \beta_\ell) + m_2(\beta'_\ell - \beta_\ell)], \\
L_{L,0}^{+\frac{1}{2}-\frac{1}{2}} &= L_{L,0}^{-\frac{1}{2}+\frac{1}{2}} = 0, \quad L_{L,0}^{-\frac{1}{2}-\frac{1}{2}} = [m_1(\beta'_\ell - \beta_\ell) + m_2(\beta'_\ell + \beta_\ell)],
\end{aligned} \tag{D13}$$

$$\begin{aligned}
L_{R,0}^{+\frac{1}{2}+\frac{1}{2}} &= -[m_1(\beta'_\ell - \beta_\ell) + m_2(\beta'_\ell + \beta_\ell)], \\
L_{R,0}^{+\frac{1}{2}-\frac{1}{2}} &= L_{R,0}^{-\frac{1}{2}+\frac{1}{2}} = 0, \quad L_{R,0}^{-\frac{1}{2}-\frac{1}{2}} = -[m_1(\beta'_\ell + \beta_\ell) + m_2(\beta'_\ell - \beta_\ell)].
\end{aligned} \tag{D14}$$

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