Testing a spin-2 mediator by angular observables in $b \rightarrow s\mu^+\mu^-$

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We consider the effects of the spin-2 particle in the $b \rightarrow s\mu^+\mu^-$ transition assuming that the spin-2 particle couples in a flavor-nonuniversal way to b and s quarks and in the leptonic sector couples only to the muons, thereby only contributing to the process $b \rightarrow s\mu^+\mu^-$. The $B_s - \bar{B}_s$ transition gives the strong constraint on the coupling of the spin-2 mediator and b and s quarks, while the observed discrepancy from the standard model prediction for the muon anomalous magnetic moment $(g-2)_\mu$ serves to constrain the μ coupling to a spin-2 particle. We find that the spin-2 particle can modify the angular observables in the $B \rightarrow K\mu^+\mu^-$ and $B \rightarrow K^*\mu^+\mu^-$ decays and produce effects that do not exist in the standard model. The generated forward-backward asymmetries in these processes can reach 15%, while other observables for these decays receive tiny effects.

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I. INTRODUCTION

Over the past few years, many theoretical and experimental studies of B -meson processes were done to test viability of the standard model (SM). On the experimental side, precision measurements of many variables have been performed, while on the theoretical side, complicated B -meson dynamics, which requires precise knowledge of many hadronic parameters, obscured precise calculations. Although the lattice QCD was able to provide us with very precise determinations of many variables, some information is still lacking.

Such precision on both sides enables us to establish a discrepancy between theoretical prediction and experimental results in the charged-current transition $b \rightarrow c\tau\nu$ as well as in the flavor-changing neutral current transition (FCNC) of $b \rightarrow s\mu^+\mu^-$. These discrepancies can be searched for in different physical observables, which are sensitive to the new physics. The LHCb Collaboration performed measurements of the relevant variables in the FCNC transition $b \rightarrow s\mu^+\mu^-$ within B -meson systems. The branching ratio of the simplest process of the type $B_s \rightarrow \mu^+\mu^-$ were measured by the both CMS and LHCb Collaborations [1], and the result was found to be somewhat lower than the predicted one [2]. The ratio of the branching fractions for the muonic and electronic mode of the $B \rightarrow K\ell^+\ell^-$ decay

in the low dilepton invariant-mass region $1 \text{ GeV}^2 \leq q^2 \leq 6 \text{ GeV}^2$ [3], known as the R_K anomaly, exhibits 2.4σ deviations below the SM prediction $R_K^{\text{SM}} = 1.00(1)$ [4]. Similarly, the LHCb Collaboration measured that in the case of $B \rightarrow K^*\ell^+\ell^-$ the experimental result for R_K^* disagrees with the SM predictions at the level of $(2.2-2.4)\sigma$ [5]. The deviations from the SM predictions triggered the questioning of the viability of lepton-flavor universality [6], and many explanations of the $R_{K^{(*)}}$ anomaly have been offered in the literature, mostly assuming some scenario of new physics (NP), e.g., Refs. [7–37].

Apart from the observables $R_{K^{(*)}}$, the angular distributions in $B \rightarrow K^{(*)}\ell^+\ell^-$ decays as a function of dilepton invariant mass squared (q^2) are also sensitive to NP effects. The kinematic distribution of the decay products of $B \rightarrow K\ell^+\ell^-$ at any particular q^2 value depends on the angle between the directions of B and of ℓ^- in the center-of-mass frame of the lepton pair. A forward-backward asymmetry constructed at this angle is zero in the SM; therefore, a nonzero measurement of the asymmetry will be a clean signal of NP [38,39]. Similarly, the full distribution of the decay $B \rightarrow K^*(K\pi)\ell^+\ell^-$ at any particular q^2 can be expressed as a triply differential cross section in three angles, namely, a) the angle between the direction of K in the K^* rest frame and K^* in the B rest frame, b) the angle between ℓ^- in the dilepton rest frame and the direction of flight of two leptons in the B rest frame, and c) the angle between $\ell^+\ell^-$ and $K\pi$ decay planes in the B rest frame. Different forward-backward or CP asymmetries can be then constructed by integrating over these angles so as to extract information about NP in detail [40].

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In this study, we investigate the effects of spin-2 particle, which contributes to the $b \rightarrow s\mu^+\mu^-$ transitions. We assume the most general coupling of such a massive spin-2 particle, allowing for the flavor violation. As suggested by Ref. [41], the effective Lagrangian study implies that the lowest-order couplings of a spin-2 state with quark fields are quite similar to the general-relativistic couplings of the graviton to the energy-momentum tensor. The couplings of the spin-2 particle with b , s and with $\mu\mu$ are highly constrained by $B_s - \bar{B}_s$ mixing and $g - 2$ of the muon. We show that this scenario therefore cannot explain the experimentally observed R_K anomalies. However, the spin-2 framework induces forward-backward asymmetries in $B \rightarrow K^{(*)}\mu^+\mu^-$, which are either zero or very small in the SM and give a clean test of our spin-2 model. We consider in Sec. II the framework of our study, introducing the interaction of spin-2 particle with fermions of the SM in $b \rightarrow s\mu^+\mu^-$, and then we constrain model parameters using the $B_s - \bar{B}_s$ oscillations and muon anomalous magnetic moment $(g - 2)_\mu$. Results of our testing of the spin-2 mediator in $B \rightarrow K^{(*)}\mu^+\mu^-$ decays are presented in Sec. III. In Sec. IV, we discuss numerical results and shortly summarize the obtained results.

II. FRAMEWORK

The spin-2 particles might appear in the particle physics due to a number of reasons, for example, in lattice QCD as spin-2 long-lived glueballs [42]. They also appear in the modified gravity. Such models contain a Kaluza Klein tower of gravitons in theories with large extra dimensions [43,44], wrapped or extra dimensions [45,46]. We follow the most general approach independent of the UV completion as given in Ref. [41]. A complex symmetric spin-2 field $G_{\mu\nu}$ having mass M can be written in the Fierz-Pauli form [47]

$$\mathcal{L}_{PF} = -\frac{1}{2}G_{\mu\nu}^\dagger(\square + M^2)G^{\mu\nu} + \frac{1}{2}G_\mu^{\mu\dagger}(\square + M^2)G_\nu^\nu - G_{\mu\nu}^\dagger\partial^\mu\partial^\nu G_\rho^\rho + G_{\mu\nu}^\dagger\partial^\mu\partial^\rho G_\rho^\nu + \text{H.c.} \quad (1)$$

Following Ref. [41], we consider an effective Lagrangian for a massive spin-2 boson without any assumption on the origin of the spin-2 field. This means that it might be a nonuniversal interacting gravitational Kaluza Klein mode, related to Randall-Sundrum theory. Another possibility is that $G_{\mu\nu}$ is a bound state of a strongly coupled theory emerging above electroweak scale. As pointed out in Refs. [41,48–51], an effective Lagrangian describing interactions with the matter fields consists of operators with various dimensions. Suppose that the scale at which the new interaction emerges is well above the electroweak scale. Then, the operators in the effective Lagrangian have to be suppressed by the appropriate dimension of this scale. The dimension-4 Lagrangian with no derivative on the fermionic fields will have a form of

$$\mathcal{L}_4 \subset -G_{\mu}^{\mu}\lambda_{ij}\bar{\psi}_i\psi_j + \text{H.c.} \quad (2)$$

Such interaction is, in its form, similar to the interaction of a scalar field, although the correlator for spin-2 particle is different from the spin-0 one. In our study, we will not consider these effects since they can be reduced to the effects of the scalar/pseudoscalar operators in the effective Lagrangian describing the $b \rightarrow s\mu^+\mu^-$ transitions as performed in Ref. [38]. According to the analysis of Ref. [52] such operators, however, cannot explain $R_{K^{(*)}}$ anomalies.

The most general coupling of the spin-2 particle with the fermions at the one-derivative order is given by the dimension-5 Lagrangian [41,48–51]

$$\mathcal{L}_5 = -\frac{i}{4\Lambda}\{a_{ij}^L[\bar{\psi}_i(\gamma_\mu\partial_\nu + \gamma_\nu\partial_\mu)P_L\psi_j] + b_{ij}^L[(\partial_\nu\bar{\psi}_i\gamma_\mu + \partial_\mu\bar{\psi}_i\gamma_\nu)P_L\psi_j]\}G^{\mu\nu} + (L \rightarrow R), \quad (3)$$

where the constants a_{ij}, b_{ij} will be taken as real and in addition we take $a_{ij} = a_{ji}, b_{ij} = b_{ji}$. Although we consider a general spin-2 particle, similarly as in Ref. [41], we will restrict the couplings by imposing

$$a_{ij}^{L,R} = -b_{ij}^{L,R}, \quad (4)$$

which appear to be valid for spin-2 particle couplings with fermions in the gravitation theory (for example, in the context of the Randall-Sundrum models with the SM fermions propagating in the bulk) [53–57]. The couplings $a_{ij}^{L,R}$ are going to be arbitrary parameters, although, due to the V-A nature of the charged weak current couplings, one expects that a_{ij}^R couplings are highly suppressed. We also consider the perturbative limits for the couplings of spin 2 with the quarks and the leptons. Since we only consider the consequences of the spin-2 field on the $b \rightarrow s\mu\mu$ transitions, we do not specify the UV completion of the theory. A detailed discussion on the range of validity of this low-energy effective theory in the context of theories with a massive general relativity in four dimensions can be found in Ref. [41].

There are also additional terms in the Lagrangian of the form

$$\mathcal{L}'_5 = -\frac{1}{2\Lambda}\eta_{\mu\nu}\{c_{ij}^L\bar{\Psi}_i i\gamma_\rho\partial^\rho\Psi_j\}G^{\mu\nu} + (L \rightarrow R), \quad (5)$$

which might contribute, too. After using the equation of motion for fermion fields, these terms will again produce scalar and pseudoscalar operator contributions being highly suppressed; see Refs. [38,51].¹

¹In Ref. [51], the authors have written the exact structure of SM fermion interactions with gravitons at the one-loop level, which could be then promoted to the most generalized interaction Lagrangian of fermions with a spin-2 field. However, since here we are interested only in leading effects, we will work only with the leading part of the Lagrangian (3) [with constraint (4) taken into account].

FIG. 1. Three-point vertex and the propagator used for our calculation, with $X_{\mu\nu}$ and $P_{\mu\nu,\rho\sigma}$ defined in Eq. (6).

Within the framework described above, the Feynman rules for the spin-2 interactions with the SM fermions in the de Donder gauge [48] are given in Fig. 1, with the symbols defined as

$$\begin{aligned}
 X_{\mu\nu} &= a_{ij}^L \gamma_\mu (k_{1\nu} + k_{2\nu}) P_L + a_{ij}^R \gamma_\mu (k_{1\nu} + k_{2\nu}) P_R \\
 P_{\mu\nu,\rho\sigma}(q^2) &= \frac{i}{q^2 - m_G^2} \left[\left(\eta_{\mu\rho} - \frac{q_\mu q_\rho}{m_G^2} \right) \left(\eta_{\nu\sigma} - \frac{q_\nu q_\sigma}{m_G^2} \right) \right. \\
 &\quad + \left(\eta_{\nu\rho} - \frac{q_\nu q_\rho}{m_G^2} \right) \left(\eta_{\mu\sigma} - \frac{q_\mu q_\sigma}{m_G^2} \right), \\
 &\quad \left. - \frac{2}{3} \left(\eta_{\mu\nu} - \frac{q_\mu q_\nu}{m_G^2} \right) \left(\eta_{\rho\sigma} - \frac{q_\rho q_\sigma}{m_G^2} \right) \right]. \quad (6)
 \end{aligned}$$

Provided the momenta transfer $q^2 \ll m_G^2$, the nonlocal interaction can be approximated to a local interaction, using Taylor's expansion

$$\begin{aligned}
 \frac{1}{q^2 - m_G^2} &= -\frac{1}{m_G^2} \left[\frac{1}{1 - q^2/m_G^2} \right] \\
 &= -\frac{1}{m_G^2} \left[1 + \frac{q^2}{m_G^2} + \frac{q^4}{m_G^4} + \dots \right]. \quad (7)
 \end{aligned}$$

On the other hand, we are interested in the most general effective Hamiltonian describing the $b \rightarrow s\mu^+\mu^-$ transitions at low energy. Such processes at scale $\mu = \mu_b = 4.8$ GeV are governed by the dimension-6 effective Hamiltonian

$$\begin{aligned}
 \mathcal{H}_{\text{eff}} &= -\frac{4G_F}{\sqrt{2}} V_{tb} V_{ts}^* \left(\sum_i C_i(\mu) \mathcal{O}_i(\mu) + C'_i(\mu) \mathcal{O}'_i(\mu) \right. \\
 &\quad \left. + \sum_{j,h.d.} C_j^{h.d.} \mathcal{O}_j^{h.d.} \right) + \text{H.c.}, \quad (8)
 \end{aligned}$$

where the operators are

$$\begin{aligned}
 \mathcal{O}_7 &= \frac{e^2}{g^2} m_b (\bar{s} \sigma_\mu P_R b) F^{\mu\nu}, \\
 \mathcal{O}'_7 &= \frac{e^2}{g^2} m_b (\bar{s} \sigma_\mu P_L b) F^{\mu\nu}, \\
 \mathcal{O}_9 &= \frac{e^2}{16\pi^2} (\bar{s} \gamma_\mu P_L b) (\bar{\ell} \gamma^\mu \ell), \\
 \mathcal{O}'_9 &= \frac{e^2}{16\pi^2} (\bar{s} \gamma_\mu P_R b) (\bar{\ell} \gamma^\mu \ell), \\
 \mathcal{O}_{10} &= \frac{e^2}{16\pi^2} (\bar{s} \gamma_\mu P_L b) (\bar{\ell} \gamma^\mu \gamma_5 \ell), \\
 \mathcal{O}'_{10} &= \frac{e^2}{16\pi^2} (\bar{s} \gamma_\mu P_R b) (\bar{\ell} \gamma^\mu \gamma_5 \ell), \\
 \mathcal{O}_S &= \frac{e^2}{16\pi^2} (\bar{s} P_R b) (\bar{\ell} \ell), \\
 \mathcal{O}'_S &= \frac{e^2}{16\pi^2} (\bar{s} P_L b) (\bar{\ell} \ell), \\
 \mathcal{O}_P &= \frac{e^2}{16\pi^2} (\bar{s} P_R b) (\bar{\ell} \gamma_5 \ell), \\
 \mathcal{O}'_P &= \frac{e^2}{16\pi^2} (\bar{s} P_L b) (\bar{\ell} \gamma_5 \ell), \\
 \mathcal{O}_T &= \frac{e^2}{16\pi^2} (\bar{s} \sigma_{\mu\nu} P_L b) (\bar{\ell} \sigma^{\mu\nu} \ell), \\
 \mathcal{O}_{T5} &= \frac{e^2}{16\pi^2} (\bar{s} \sigma_{\mu\nu} P_L b) (\bar{\ell} \sigma^{\mu\nu} \gamma_5 \ell). \quad (9)
 \end{aligned}$$

Note that the operators $\mathcal{O}_i$, $i = 7, 9, 10$, appear in the SM physics. Usually, NP also generates operators of dimension 6, although it is sometimes necessary to enlarge the base of the operator to the higher dimensions, like, for example, in the case of the two Higgs doublet model [26,27] in which the operator basis has to be enlarged by the dimension 7 operators. In the next section, we will show in our model the enlargement of the operator basis by operators of dimension 8 induced by the spin-2 particle mediation.

A. Spin-2 mediator in the $b \rightarrow s\mu^+\mu^-$ transition

We first consider the contributions of the spin-2 particle to the $b \rightarrow s\mu^+\mu^-$ amplitude, as presented in Fig. 2. By such a choice, we assume that this fermion spin-2 interaction is

FIG. 2. Feynman diagrams for the $b \rightarrow s\mu^+\mu^-$ transition with the spin-2 mediation.

flavor nonuniversal and that only relevant couplings are bs to the spin-2 particle and $\mu\mu$ to spin-2 particle couplings. In the calculation of the $b(p_b) \rightarrow s(p_s)\mu^-(k_-)\mu^+(k_+)$ amplitude, we use the following notation: $q = p_b - p_s = k_+ + k_-$, $p = p_b + p_s$, and $k = k_+ - k_-$.

The amplitude for the $b \rightarrow s\mu^+\mu^-$ transition is then of the form

$$\begin{aligned} i\mathcal{M} &= \left(\frac{-i}{4\Lambda}\right)^2 a_{sb}^i a_{\mu\mu}^j \{\bar{u}_s(p_s)[\gamma^\mu(p_s^\nu + p_b^\nu) \\ &\quad + \gamma^\nu(p_s^\mu + p_b^\mu)]P_i u_b(p_b)\} P_{\mu\nu,\rho\sigma} \\ &\quad \times \{\bar{u}_\ell(k_-)[\gamma^\rho(k_-^\sigma - k_+^\sigma) + \gamma^\sigma(k_-^\rho - k_+^\rho)]P_j v_\ell(k_+)\} \\ &= i \frac{a_{sb}^i a_{\mu\mu}^j}{16\Lambda^2 m_G^2} \{\bar{u}_s(p_s)[\gamma^\mu(p_s^\nu + p_b^\nu) \\ &\quad + \gamma^\nu(p_s^\mu + p_b^\mu)]P_i u_b(p_b)\} T_{\mu\nu}^j, \end{aligned} \quad (10)$$

where $T_{\mu\nu}$ is obtained through the contraction of the propagator with the leptonic current,

$$\begin{aligned} T_{\mu\nu}^L &= -\frac{4}{3} m_\ell [\bar{u}_\ell(k_-)P_L v_\ell(k_+) + \bar{u}_\ell(k_-)P_R v_\ell(k_+)]g_{\mu\nu} \\ &\quad + 2(k_{-\mu} - k_{+\mu})[\bar{u}_\ell(k_-)\gamma_\nu P_L v_\ell(k_+)] \\ &\quad + 2(k_{-\nu} - k_{+\nu})[\bar{u}_\ell(k_-)\gamma_\mu P_L v_\ell(k_+)], \\ T_{\mu\nu}^R &= T_{\mu\nu}^L \quad (\text{with } L \leftrightarrow R). \end{aligned} \quad (11)$$

After multiple uses of the equations of motion, we determine contributions to the scalar operators of dimension 8,

$$\begin{aligned} \mathcal{O}_S^{(q,8)} &= m_\ell m_q \frac{e^2}{16\pi^2} (\bar{s}P_R b)(\bar{\mu}\mu), \\ \mathcal{O}_{S'}^{(q,8)} &= m_\ell m_q \frac{e^2}{16\pi^2} (\bar{s}P_L b)(\bar{\mu}\mu), \end{aligned} \quad (12)$$

with $q = s, b$. The product $m_\ell m_s$, being small, can be safely neglected. We therefore consider the operators $\mathcal{O}_S^{(b,8)}, \mathcal{O}_{S'}^{(b,8)}$, henceforth referred to as $\mathcal{O}_S^{(8)}, \mathcal{O}_{S'}^{(8)}$. The appropriate Wilson coefficients for the operators proportional to $m_\ell m_b$ are

$$\begin{aligned} C_S^{(8)} &= \frac{4}{3} C_G a_{sb}^L [a_{\mu\mu}^L + a_{\mu\mu}^R], \\ C_{S'}^{(8)} &= \frac{4}{3} C_G a_{sb}^R [a_{\mu\mu}^L + a_{\mu\mu}^R], \end{aligned} \quad (13)$$

where we used

$$C_G = \frac{16\pi^2 v^2}{e^2} \frac{1}{V_{tb} V_{ts}^*} \frac{1}{16\Lambda^2 m_G^2}. \quad (14)$$

Here, the dimension-8 operators are proportional to the product of $m_b m_\mu$ with the dimension-6 operators $\mathcal{O}_S$ and $\mathcal{O}_{S'}$. However, there are new additional operators of dimension 8,

$$\begin{aligned} \mathcal{O}_i^{(8)} &= (\bar{s}\gamma^\mu i\overleftrightarrow{\partial}^\nu P_i b)(\bar{\ell}\gamma_\mu i\overleftrightarrow{\partial}_\nu \ell), \\ \mathcal{O}_{i5}^{(8)} &= (\bar{s}\gamma^\mu i\overleftrightarrow{\partial}^\nu P_i b)(\bar{\ell}\gamma_\mu i\overleftrightarrow{\partial}_\nu \gamma_5 \ell), \\ \mathcal{O}_{ij}^{(8)} &= (\bar{s}\gamma^\nu i\overleftrightarrow{\partial}^\mu P_i b)(\bar{\ell}\gamma_\mu i\overleftrightarrow{\partial}^\nu P_j \ell), \end{aligned} \quad (15)$$

with the Wilson coefficients

$$\begin{aligned} C_i^{(8)} &= -C_G a_{sb}^i (a_{\mu\mu}^L + a_{\mu\mu}^R), \\ C_{i5}^{(8)} &= -C_G a_{sb}^i (a_{\mu\mu}^R - a_{\mu\mu}^L), \\ C_{ij}^{(8)} &= -2C_G a_{sb}^i a_{\mu\mu}^j. \end{aligned} \quad (16)$$

By using the Fierz rearrangement and by applying the field equations, one can show that

$$\begin{aligned} &\int d^4x \langle l^+(k_+) l^-(k_-) s(p_s) | \mathcal{O}_i^{(8)}(x) | b(p_b) \rangle \\ &\quad \rightarrow (k \cdot p) [\bar{u}_s(p_s) \gamma^\mu P_i u_b(p_b)] [\bar{u}_\ell(k_-) \gamma_\mu v_\ell(k_+)] \\ &\int d^4x \langle l^+(k_+) l^-(k_-) s(p_s) | \mathcal{O}_{i5}^{(8)}(x) | b(p_b) \rangle \\ &\quad \rightarrow (k \cdot p) [\bar{u}_s(p_s) \gamma^\mu P_i u_b(p_b)] [\bar{u}_\ell(k_-) \gamma_\mu \gamma_5 v_\ell(k_+)] \\ &\int d^4x \langle l^+(k_+) l^-(k_-) s(p_s) | \mathcal{O}_{ij}^{(8)}(x) | b(p_b) \rangle \\ &\quad \rightarrow (k^\mu p^\nu) [\bar{u}_s(p_s) \gamma_\mu P_i u_b(p_b)] [\bar{u}_\ell(k_-) \gamma_\nu P_j v_\ell(k_+)], \end{aligned} \quad (17)$$

from which we can see that effectively the dimension-8 operators $\mathcal{O}_i^{(8)}$ and $\mathcal{O}_{i5}^{(8)}$ have the same Lorentz structure as the dimension-6 operators, the matrix elements of which can be calculated. Namely, we notice that $\mathcal{O}_i^{(8)} \sim \mathcal{O}_9, \mathcal{O}'_9$ and $\mathcal{O}_{i5}^{(8)} \sim \mathcal{O}_{10}, \mathcal{O}'_{10}$. The results of the global fit analysis [52] suggest that the B-physics anomalies can be explained by the NP for the Wilson coefficients which satisfy the relation $C_9 \simeq -C_{10}$. This indicates that the new interactions should be of V-A type. To follow such a V-A frame of NP, we impose the condition

$$a_{sb}^R, a_{\mu\mu}^R = 0 \rightarrow C_i^{(8)} = -C_{i5}^{(8)}. \quad (18)$$

The operator $\mathcal{O}_{LL}^{(8)}$, as can be seen from (17), is a genuine operator of dimension 8 that cannot be reduced to the standard dimension-6 operators. This means that we have three independent operators of the type $\mathcal{O}_L^{(8)}, \mathcal{O}_{L5}^{(8)}$, and $\mathcal{O}_{LL}^{(8)}$ and two scalar operators $\mathcal{O}_S^{(b,8)} \equiv \mathcal{O}_S^{(8)}$ and $\mathcal{O}_{S'}^{(b,8)} \equiv \mathcal{O}_{S'}^{(8)}$ [Eq. (12)] with the appropriate Wilson coefficients:

FIG. 3. A tree-level contribution to $B_s - \bar{B}_s$ mixing.

$$\begin{aligned}
 C_L^{(8)} &= -C_G a_{sb}^L a_{\mu\mu}^L, \\
 C_{LS}^{(8)} &= C_G a_{sb}^L a_{\mu\mu}^L, \\
 C_{LL}^{(8)} &= -2C_G a_{sb}^L a_{\mu\mu}^L, \\
 C_S^{(8)} &= \frac{4}{3} C_G a_{sb}^L a_{\mu\mu}^L, \\
 C_{S'}^{(8)} &= 0.
 \end{aligned} \tag{19}$$

In the following, we will consider effects of this operator basis. The new contribution can appear in the decay amplitudes for $B_s \rightarrow \mu^+ \mu^-$, $B \rightarrow K \mu^+ \mu^-$, and $B \rightarrow K^* \mu^+ \mu^-$.

B. Constraints on spin-2 couplings

Our model-independent analysis aimed at determining spin-2 mediator effects in $b \rightarrow s \mu^+ \mu^-$ requires knowledge of the bs -spin 2 coupling a_{sb}^L and $\mu\mu$ -spin 2 coupling $a_{\mu\mu}^L$. First, we discuss the determination of a_{sb}^L .

C. $B_s - \bar{B}_s$ mixing

The coupling of the b and s quarks with the spin-2 particle can be extracted from the $B_s - \bar{B}_s$ mixing observables. The spin-2 tree-level contribution to the $B_s - \bar{B}_s$ transition is presented in Fig. 3. The most general effective Hamiltonian of dimension 6 for the $\Delta B = 2$ process [58–62] can be written as

$$H_{\text{eff}}^{\Delta B=2} = \sum_{i=1}^5 C_i Q_i + \sum_{i=1}^3 \tilde{C}_i \tilde{Q}_i. \tag{20}$$

The full set of operators Q_i and $\tilde{Q}_i$ is presented in Refs. [58–63]. The current experimental result for $\Delta m_{B_s} = (1.1688 \pm 0.0014) \times 10^{-11}$ GeV [64] is consistent with the SM result $(1.10083_{-0.038}^{+0.054}) \times 10^{-11}$ GeV [65].

In the case of the spin-2 mediator, the operators are of dimension 8. These operators can be written as the product of quark masses and the dimension-6 operators $Q_i, \tilde{Q}_i$, with $i = 2, 3, 4, 5$,

$$\begin{aligned}
 Q_2^{(8)} &= m_s^2 (\bar{s}_R^\alpha b_L^\alpha) (\bar{s}_R^\beta b_L^\beta), & \tilde{Q}_2^{(8)} &= m_b^2 (\bar{s}_L^\alpha b_R^\alpha) (\bar{s}_L^\beta b_R^\beta), \\
 Q_3^{(8)} &= m_s^2 (\bar{s}_R^\alpha b_L^\beta) (\bar{s}_R^\beta b_L^\alpha), & \tilde{Q}_3^{(8)} &= m_b^2 (\bar{s}_L^\alpha b_R^\beta) (\bar{s}_L^\beta b_R^\alpha), \\
 Q_4^{(8)} &= m_b m_s (\bar{s}_R^\alpha b_L^\alpha) (\bar{s}_L^\beta b_R^\beta), & Q_5^{(8)} &= m_b m_s (\bar{s}_R^\alpha b_L^\beta) (\bar{s}_L^\beta b_R^\alpha).
 \end{aligned} \tag{21}$$

There are also two additional operators:

$$\begin{aligned}
 Q_6^{(8)} &= (\bar{s}^\alpha \gamma^\mu i \overleftrightarrow{\partial}_\nu P_L b^\alpha) (\bar{s}^\beta \gamma^\mu i \overleftrightarrow{\partial}_\nu P_L b^\beta), \\
 Q_7^{(8)} &= (\bar{s}^\alpha \gamma^\nu i \overleftrightarrow{\partial}_\mu P_L b^\alpha) (\bar{s}^\beta \gamma^\mu i \overleftrightarrow{\partial}_\nu P_L b^\beta).
 \end{aligned} \tag{22}$$

By the use of the equations of motion, the operator $Q_6^{(8)}$ can be reduced to $m_b^2 Q_1$ with $Q_1 = (\bar{s}_L^\alpha \gamma^\mu b_L^\alpha) (\bar{s}_L^\beta \gamma^\mu b_L^\beta)$. In our numerical study, the operators proportional to m_s^2 and $m_b m_s$ can be safely neglected. The corresponding Wilson coefficients of the operators are

$$\begin{aligned}
 C_2^{(8)} &= -C_3^{(8)} = \frac{1}{2} C_4^{(8)} = \tilde{C}_2^{(8)} = -\tilde{C}_3^{(8)} \\
 &= \frac{1}{2} C_5^{(8)} = -\frac{(a_{sb}^L)^2}{\Lambda^2 m_G^2} \left(\frac{8}{3} \right), \\
 C_6^{(8)} &= C_7^{(8)} = 4 \frac{(a_{sb}^L)^2}{16 \Lambda^2 m_G^2}.
 \end{aligned} \tag{23}$$

The hadronic matrix elements of the operators $m_b^2 (Q_1, \tilde{Q}_{2,3})$ between the neutral B_s mesons and the corresponding Wilson coefficients at the scale Λ evolved to the hadronic scale with the anomalous dimension matrices can be found in Ref. [58]. We added the contribution of the spin-2 mediator to the mass difference of the $B_s - \bar{B}_s$ system to the SM contribution and write it schematically as

$$\Delta M_s = (\Delta M_s)_{\text{SM}} + (\Delta M_s)_{m_b^2 Q_1} + (\Delta M_s)_{m_b^2 \tilde{Q}_2} + (\Delta M_s)_{m_b^2 \tilde{Q}_3}. \tag{24}$$

The results with spin-2 contributions are given in Appendix A. Note that the spin-2 contributions contain

FIG. 4. Combined bound from $B_s - \bar{B}_s$ mixing and the top decay width on a_{sb}^L as a function of m_G .

the common proportionality factor $\frac{(a_{sb}^L)^2}{\Lambda^2 m_G^2}$. Following the calculation of Ref. [58], we determine bound on the bs -spin 2 particle coupling:

$$\left(\frac{a_{sb}^L}{\Lambda m_G}\right)^2 < 1.4 \times 10^{-12} \text{ GeV}^{-4}. \quad (25)$$

One can find straightforwardly that, for $\Lambda \simeq 1 \text{ TeV}$ and $m_G \simeq 500 \text{ GeV}$, $a_{sb}^L \approx 0.5$. We point out that, due to the $SU(2)_L$ symmetry, one can relate the coupling of the top quark, charm quark, and spin-2 particle to the coupling of bottom quarks, strange quarks, and spin-2 mediator. The top-quark decay channel $t \rightarrow c G_{\mu\nu}$ is allowed when the mass of the spin-2 particle is less than the top-quark mass. Using $a_{sb}^L \simeq a_{tc}^L$, the decay width of this mode is given by

$$\Gamma_t^{\text{NP}} = \frac{|a_{sb}^L|^2 m_t^7}{192\pi \Lambda^2 m_G^4} \left(1 - \frac{m_G^2}{m_t^2}\right)^4 \left(2 + 3 \frac{m_G^2}{m_t^2}\right). \quad (26)$$

The SM decay width for $m_t = 173.3 \text{ GeV}$ is 1.35 GeV , whereas the measured decay width is $1.41_{-0.15}^{+0.19} \text{ GeV}$. The top decay width might give a strong bound on a_{sb}^L for $m_G < m_t$. We show in Fig. 4 the combined bound from the $B_s - \bar{B}_s$ mixing (25) and the top decay width (26) on a_{sb}^L as a function of m_G for $\Lambda = 1000 \text{ GeV}$. We consider $\Lambda \simeq 1000 \text{ GeV}$ along with m_G in the mass range 10–150 GeV so that the current LHC data do not constrain the parameter space considered here. In the case of masses above 150 GeV, the decay channel to massive gauge bosons or to a pair of Higgs bosons opens up. These channels have been explored in the LHC, and limits are obtained on the parameter space of the respective models. Similarly, in the case of spin-2 particle mass below 10 GeV, the upson decay to a dimuon final state will open, and this process has been studied in detail. This in turn leads to strong constraints on spin-2 mass below 10 GeV.

FIG. 5. The lowest-order contribution of the spin-2 particle.

D. Constraints from $(g-2)_\mu$

The experimental and theoretical results for the muon anomalous magnetic moment $a_\mu = 1/2(g-2)$ differ. The recent update of the well-established difference is $\Delta a_\mu = a_\mu^{\text{exp}} - a_\mu^{\text{th}} = (3.11 \pm 0.77) \times 10^{-9}$ [66]. In the presence of the spin-2 particle, the lowest-order contribution to the muon anomalous magnetic moment is shown in Fig. 5. We calculate the spin-2 mediator contributions to the muon anomalous magnetic moment $(g-2)_\mu$ and find

$$\begin{aligned} & \left(\frac{a_{\mu\mu}^L}{4\pi}\right)^2 \frac{2m_\mu}{(4\Lambda)^2} \int_0^1 dx dy dz \delta(x+y+z-1) \bar{u}(p') \\ & \times \left[i \frac{\sigma^{\mu\nu} q_\nu}{2m_\mu} \left(-\frac{2m_\mu^3}{3\Delta} (6z^4 + 26z^3 \right. \right. \\ & \left. \left. + 29z^2 - 5z - 11) - \frac{4m_\mu}{3} \left\{ \log \left[\frac{\Lambda_N^2 + \Delta}{\Delta} \right] \right. \right. \right. \\ & \left. \left. \left. + \frac{\Delta}{\Lambda_N^2 + \Delta} - 1 \right\} (6z^2 + 14z + 13) \right) \right] u(p), \quad (27) \end{aligned}$$

where $\Delta = (1-z)^2 m_\mu^2 + z m_G^2$. The diagram diverges, and in principle, one has to use the full theory containing the spin-2 particle to calculate it. To estimate the size of the $a_{\mu\mu}^L/\Lambda$ coupling, we assume cutoff regularization and set the cutoff scale to be almost equal to the scale of spin-2 theory $\Lambda_N \simeq \Lambda$. The lower bound obtained on $a_{\mu\mu}^L/\Lambda$ from $(g-2)_\mu$ as a function of m_G is presented in Fig. 6, which will be used in our further calculations. We have tried to constrain $a_{\mu\mu}^L$ also by considering the decay channel $Z \rightarrow 4\mu$. This decay channel is studied at the LHC, and an upper bound is obtained for the total decay width. The presence of the spin-2 particle gives an additional contribution to this process through the $a_{\mu\mu}^L$ coupling. We checked that the constraint obtained for $a_{\mu\mu}^L$ from the Z decay width to four muons is much weaker than the constraints coming from the muon anomalous magnetic moment. In addition to the constraints discussed above, we

FIG. 6. The bound on $a_{\mu\mu}^L/\Lambda$ as a function of m_G from the muon anomalous magnetic moment.

also consider the dilepton invariant-mass distribution in $pp \rightarrow \mu^+\mu^-$ at the 13 TeV LHC with 36.1 fb^{-1} of data [67]. We have performed our computation in MADGRAPH [68] and find that for the parameter space considered in our model with $\Lambda \simeq 1000 \text{ GeV}$ the LHC data in the low-mass range does not constrain the scenario considered here. The result remains the same even if we allow the spin-2 mediator to couple to all the quarks with the coupling of the same order as a_{sb}^L . This is mainly because the couplings of the spin-2 particle with quarks and leptons are suppressed by Λ as opposed to the Z' models that are currently being explored and constrained by the LHC dilepton data [69].

III. EFFECT OF SPIN 2 IN EXCLUSIVE DECAYS $B \rightarrow K^{(*)}\mu^+\mu^-$

The $b \rightarrow s\mu^+\mu^-$ transition can occur in a number of exclusive decay channels. The leptonic decay mode $B_s \rightarrow \mu^-\mu^-$ might also receive contributions from the NP. The $B_s \rightarrow \mu^-\mu^-$ amplitude arises from the contributions of axial,

scalar, and pseudoscalar lepton currents as discussed in Refs. [38,39]. In the case of the spin-2 mediator, all relevant contributions are either proportional to the muon mass, leading to the negligible effects, or kinematically equal to zero.

We will show that higher-dimensional operators (16) will produce new angular dependencies in exclusive $B \rightarrow K^{(*)}\mu^+\mu^-$ decays. The opportunity to use the measurement of higher moments of the differential angular distributions to test physics beyond the SM coming from contributions of $\mathcal{O}(m_B^2/\Lambda^2)$ suppressed higher-dimensional operators and then distinguishing these contributions from the estimated higher-order QED corrections was discussed in [70], where the authors have used the general helicity formalism to derive full angular distributions for the complete dimension-six effective Hamiltonian in those decays.

A. Spin-2 mediator in $B \rightarrow K\mu^+\mu^-$

In our study, we follow the approach and notations introduced in Refs. [38,39] (for details, see Appendix B). The angular observables are very useful for identifying NP effects in $B \rightarrow K\mu^+\mu^-$. The angle θ is defined as the angle between the direction of B and the lepton ℓ^- in the center-of-mass frame of the lepton pair, and q^2 is running from $q_{\min}^2 = 4m_\ell^2$ to $q_{\max}^2 = (m_B - m_K)^2$. As usual, we follow definitions $\beta_\ell(q^2) = \sqrt{1 - 4m_\ell^2/q^2}$ and $\lambda(q^2) = q^4 + m_B^4 + m_K^4 - 2(m_B^2m_K^2 + m_B^2q^2 + m_K^2q^2)$. Using this notation, we write down the double differential decay rate with respect to q^2 and $\cos\theta$ for the process $B \rightarrow K\ell^+\ell^-$ as

$$\frac{d^2\Gamma_\ell}{dq^2 d\cos\theta} = a_\ell(q^2) + \tilde{b}_\ell(q^2)\cos\theta + c_\ell(q^2)\cos^2\theta + \tilde{d}_\ell(q^2)\cos^3\theta + \tilde{e}_\ell(q^2)\cos^4\theta, \quad (28)$$

where we have

$$\begin{aligned} a_\ell(q^2) &= \mathcal{N}(q^2) \left[q^2(\beta_\ell^2(q^2)|F_S^{\text{NP}}(q^2)|^2 + |F_P(q^2)|^2) + \frac{\lambda(q^2)}{4}(|F_A(q^2)|^2 + |F_V(q^2)|^2) \right. \\ &\quad \left. + 4m_\ell^2m_B^2|F_A(q^2)|^2 + 2m_\ell(m_B^2 - m_K^2 + q^2)\text{Re}(F_P(q^2)F_A^*(q^2)) \right], \\ \tilde{b}_\ell(q^2) &= 2\mathcal{N}(q^2) \left[m_\ell\sqrt{\lambda(q^2)}\beta_\ell(q^2)\text{Re}(F_S^{\text{NP}}(q^2)F_V^*(q^2)) + q^2\text{Re}(F_P^{\text{NP}}(q^2)F_P^*(q^2)) + \frac{\lambda(q^2)}{4}\text{Re}(F_A^{\text{NP}}(q^2)F_A^*(q^2)) \right. \\ &\quad \left. + F_V^{\text{NP}}(q^2)F_V^*(q^2) + 4m_\ell^2m_B^2\text{Re}(F_A^{\text{NP}}(q^2)F_A^*(q^2)) + m_\ell(m_B^2 - m_K^2 + q^2)\text{Re}(F_P^{\text{NP}}(q^2)F_A^*(q^2) + F_P(q^2)F_A^{\text{NP}*}(q^2)) \right. \\ &\quad \left. + \frac{\sqrt{\lambda(q^2)}}{8}\beta_\ell(q^2)f_+(q^2)C_{LL}^{(8)} \left\{ 4\frac{m_\ell^2}{q^2}C_{10}((m_B^2 - m_K^2)^2f_0(q^2) - \lambda(q^2)f_+(q^2)) - \lambda(q^2)(F_V(q^2) - F_A(q^2)) \right\} \right], \\ c_\ell(q^2) &= \mathcal{N}(q^2) \left[-\frac{\lambda(q^2)}{4}\beta_\ell^2(q^2)(|F_A(q^2)|^2 + |F_V(q^2)|^2) + q^2|F_P^{\text{NP}}(q^2)|^2 + \frac{\lambda(q^2)}{4}(|F_A^{\text{NP}}(q^2)|^2 \right. \\ &\quad \left. + |F_V^{\text{NP}}(q^2)|^2) + 4m_\ell^2m_B^2|F_A^{\text{NP}}(q^2)|^2 + 2m_\ell(m_B^2 - m_K^2 + q^2)\text{Re}(F_P^{\text{NP}}(q^2)F_A^{\text{NP}*}(q^2)) \right. \\ &\quad \left. + \frac{\beta_\ell^2(q^2)}{8}f_+(q^2)\lambda(q^2) \left\{ C_{LL}^{(8)}(C_{LL}^{(8)}f_+(q^2)(\lambda(q^2) + 2m_\ell^2(2(m_B^2 + m_K^2) - q^2)) + \frac{8}{m_b}F_S^{\text{NP}}(q^2)) \right\} \right], \end{aligned}$$

$$\begin{aligned}
\tilde{d}_\ell(q^2) &= \mathcal{N}(q^2) \left[\frac{1}{4} \beta_\ell^3(q^2) \lambda^{3/2}(q^2) f_+(q^2) C_{LL}^{(8)} (F_V(q^2) - F_A(q^2)) \right. \\
&\quad \left. - \frac{\lambda(q^2)}{4} \beta_\ell^2(q^2) 2\text{Re}(F_A^{\text{NP}}(q^2) F_A^*(q^2) + F_V^{\text{NP}}(q^2) F_V^*(q^2)) \right], \\
\tilde{e}_\ell(q^2) &= \mathcal{N}(q^2) \left[-\frac{\lambda(q^2)}{4} \beta_\ell^2(q^2) (|F_A^{\text{NP}}(q^2)|^2 + |F_V^{\text{NP}}(q^2)|^2) - \frac{\beta_\ell^4}{8} f_+(q^2) \lambda(q^2)^2 (C_{LL}^{(8)})^2 \right]
\end{aligned} \tag{29}$$

with

$$\mathcal{N}(q^2) = \frac{G_F^2 \alpha^2 |V_{tb} V_{ts}^*|^2}{512 \pi^5 m_B^3} \beta_\ell(q^2) \sqrt{\lambda(q^2)}. \tag{30}$$

Functions $\tilde{b}(q^2)$, $\tilde{d}(q^2)$, and $\tilde{e}(q^2)$ are zero in the SM. The terms proportional to $\cos \theta$ might appear in some extensions of the SM (see, e.g., Ref. [38]), while the terms $\tilde{d}(q^2)$ and $\tilde{e}(q^2)$ are a genuine consequence of the spin-2 mediator.

The SM form factors are extended by the effects of the spin-2 particle. Following the notation of Ref. [38], we write

$$\begin{aligned}
F_V(q^2) &= C_9 f_+(q^2) + \frac{2m_b}{m_B + m_K} C_7 f_T(q^2), \\
F_V^{\text{NP}}(q^2) &= -\beta_\ell \sqrt{\lambda(q^2)} C_L^{(8)} f_+(q^2), \\
F_A(q^2) &= C_{10} f_+(q^2), \\
F_A^{\text{NP}}(q^2) &= -\beta_\ell \sqrt{\lambda(q^2)} C_{L5}^{(8)} f_+(q^2), \\
F_P(q^2) &= -m_\ell C_{10} \left[f_+(q^2) - \frac{m_B^2 - m_K^2}{q^2} (f_0(q^2) - f_+(q^2)) \right], \\
F_P^{\text{NP}}(q^2) &= m_\ell \beta_\ell \sqrt{\lambda(q^2)} C_{L5}^{(8)} \left[f_+(q^2) - \frac{m_B^2 - m_K^2}{q^2} (f_0(q^2) - f_+(q^2)) \right], \\
F_S^{\text{NP}}(q^2) &= m_\ell \frac{m_B^2 - m_K^2}{2} (C_S^{(8)} + C_{S'}^{(8)}) f_0(q^2).
\end{aligned} \tag{31}$$

The spin-2 particle contributions to the total decay width are proportional to $(C_i^{(8)})^2 \sim 1/(\Lambda^4 m_G^4)$, and therefore they are insignificant.

The interference terms proportional to $\cos \theta$ and $\cos^3 \theta$ will, however, contribute to the forward-backward asymmetry (A_{FB}), which is zero in the SM. The forward-backward asymmetry is defined as

$$A_{\text{FB}} = \frac{1}{\Gamma_\ell} \int dq^2 \left(\int_0^1 - \int_{-1}^0 \right) d \cos \theta \left(\frac{d^2 \Gamma_\ell}{dq^2 d \cos \theta} \right) = \frac{1}{\Gamma_\ell} \int dq^2 \left(\tilde{b}_\ell(q^2) + \frac{\tilde{d}_\ell(q^2)}{2} \right), \tag{32}$$

where the total decay rate is given by

$$\Gamma_\ell = 2 \int dq^2 \left(a_\ell(q^2) + \frac{c_\ell(q^2)}{3} + \frac{\tilde{e}_\ell(q^2)}{5} \right). \tag{33}$$

In our calculation, we consider only the region of the square of the invariant dilepton mass $1 \leq q^2 \leq 6 \text{ GeV}^2$. The odd power of $\cos \theta$ proportional to $\tilde{b}_\ell(q^2)$ and $\tilde{d}_\ell(q^2)$, which contains the interference term of the SM with the spin-2 contribution, survives. A nonzero measurement of this asymmetry therefore will be a clear signal of this NP scenario. It is interesting that spin-2 dimension-8 operators

introduce dependence on $\cos \theta$ and $\cos^3 \theta$, not present in the SM. The terms $a_\ell(q^2)$, $c_\ell(q^2)$, and $\tilde{e}_\ell(q^2)$ contributing to the total decay width are sensitive to the spin-2 contribution in the form of $\text{NP}^2 \approx 1/(16\Lambda^2 m_G^2)^2$. This in turn leads to the increase of the order of $\sim 1\text{--}2\%$ as presented in Fig. 7. Our hypothesis on the muon specific interaction of the spin-2 particle implies that the first lepton generation does not receive the spin-2 corrections. However, due to the tiny increase of the differential branching ratio (BR) $\text{BR}(K \rightarrow K \mu^+ \mu^-)$, the ratio R_K is only insignificantly different from 1. Obviously, the interaction of the spin-2 mediator with the b and s quarks as well as muons, without additional new particles, like

FIG. 7. $\text{Br}(B \rightarrow K \mu^+ \mu^-)$ as a function of spin-2 mass m_G for $\Lambda = 1000$ GeV in the presence of the spin-2 contribution (red), taking into account the constraints from the muon anomalous magnetic moment, $B_s - \bar{B}_s$ mixing, and the top decay width. The SM value for the BR is shown by a blue line.

FIG. 8. The forward-backward asymmetry A_{FB} as a function of spin-2 mass m_G , taking into account the constraints from the $B_s - \bar{B}_s$ mixing, the top decay width, and muon anomalous magnetic moment for $\Lambda = 1000$ GeV in the presence of spin-2 contribution. Note that the SM value of A_{FB} is zero.

vectorlike fermions or gauge bosons, cannot explain the $R_{K^{(*)}}$ anomaly. We present in Fig. 8 A_{FB} as a function of m_G , for all the points satisfying the bounds coming from the $B_s - \bar{B}_s$ mixing, the top decay width, and the muon anomalous magnetic moment. For the rather low values of the spin-2 particle mass $m_G \simeq 20\text{--}40$ GeV, the generated forward-backward asymmetry can be as large as 20%.

There is $\cos^4 \theta$ dependence as a spin-2 particle effect in Eq. (28), proportional to $\tilde{e}_\ell(q^2)$. We find that $E_\ell = (1/\Gamma_\ell) \int dq^2 \tilde{e}_\ell(q^2)$, being a contribution of the order $(1/\Lambda)^4$, is insignificant and can be only $\sim 1\%$ of the leading contribution coming from the term proportional to $a_\ell(q^2)$, for m_G in the range $20\text{--}50$ GeV and $\Lambda = 1000$ GeV. Another observable considered in Refs. [39,71] is $F_H^\ell = (2/\Gamma_\ell) \int dq^2 (a_\ell(q^2) + c_\ell(q^2))$. This term, being proportional to m_μ^2 , is small in the SM. The inclusion of the spin-2 contributions modifies negligibly this observable by the terms proportional to $(1/\Lambda)^4$.

It is worth pointing out that the left-handed nature of the muonic coupling to the spin-2 particle will induce couplings of the spin-2 particle to neutrinos. That will imply that the branching ratio of $B \rightarrow K^{(*)} \nu \nu$ might obtain

additional contributions. The Belle Collaboration determined the upper bounds for the branching fractions of these decays [72]. The effect of the spin-2 particle in $\text{Br}(B \rightarrow K^{(*)} \nu \bar{\nu})$ is almost equal to the effect of spin 2 in the branching ratio for $B \rightarrow K^{(*)} \mu^+ \mu^-$. Therefore, it contributes insignificantly to the SM prediction.

B. Spin-2 mediator in the $B \rightarrow K^*(\rightarrow K\pi) \mu^+ \mu^-$ decay

The spin-2 particle can also affect the process $B \rightarrow K^*(\rightarrow K\pi) \mu^+ \mu^-$. The relevant form factors for the $B \rightarrow K^*$ matrix elements are defined in Refs. [40,73] (for details, see Appendix B), with $q^\mu = p^\mu - p_{K^*}^\mu$ and ϵ_μ as the polarization vector of K^* . The angular distribution of the four-body decay $B \rightarrow K^*(\rightarrow K\pi) \ell^+ \ell^-$ has been discussed extensively in the literature (see, e.g., Refs. [40,73]). We determine contributions of the operators $\mathcal{O}_L^{(8)}$, $\mathcal{O}_{LS}^{(8)}$, and $\mathcal{O}_{LL}^{(8)}$ to the $B \rightarrow K^* \mu^+ \mu^-$ decay amplitude

$$\frac{d^4 \Gamma}{dq^2 d \cos \theta_\ell d \cos \theta_{K^*} d \phi} = \frac{9}{32\pi} I(q^2, \theta_\ell, \theta_{K^*}, \phi), \quad (34)$$

where

$$\begin{aligned} I(q^2, \theta_\ell, \theta_{K^*}, \phi) = & I_1^s \sin^2 \theta_{K^*} + I_1^c \cos^2 \theta_{K^*} + (I_2^s \sin^2 \theta_{K^*} + I_2^c \cos^2 \theta_{K^*}) \cos 2\theta_\ell, \\ & + I_3 \sin^2 \theta_{K^*} \sin^2 \theta_\ell \cos 2\phi + I_4 \sin 2\theta_{K^*} \sin 2\theta_\ell \cos \phi + I_5 \sin 2\theta_{K^*} \sin \theta_\ell \cos \phi \\ & + (I_6^s \sin^2 \theta_{K^*} + I_6^c \cos^2 \theta_{K^*}) \cos \theta_\ell + I_7 \sin 2\theta_{K^*} \sin \theta_\ell \sin \phi \\ & + I_8 \sin 2\theta_{K^*} \sin 2\theta_\ell \sin \phi + I_9 \sin^2 \theta_{K^*} \sin^2 \theta_\ell \sin 2\phi. \end{aligned} \quad (35)$$

The angular coefficients I_i 's in the presence of standard dimension-6 operators are dependent on q^2 . The full expressions for I_i as a function of the transversity amplitudes are given in Appendix B. We define θ_{K^*} as the angle of the direction of flight of K in the K^* rest frame with respect to the direction of flight of K^* in the B rest frame and θ_ℓ as the angle of the direction of flight of ℓ^- in the dilepton rest frame with respect to the direction of flight of the two leptons in the B rest frame. The angle between the plane formed by $K\pi$ and $\ell\ell$ is denoted by ϕ .

The angular coefficients are combinations of the transversity amplitudes $A_{\perp,\parallel,0,t}^{L,R}$, which are listed in Ref. [40]. We have also followed the convention of Ref. [40] for the polarization vector of the virtual gauge boson $\epsilon_{V^*}^\mu(n)$, $n = \pm$ (transverse), $n = 0$ (longitudinal), or $n = t$ (timelike) and the K^* polarization vector $\epsilon_{K^*}^\mu(m)$, with $m = \pm, 0$. The transversity amplitudes $A_{\perp,\parallel,0,t}^{L,R}$ in our case get modified by the dimension-8 operators and are listed below,

$$\begin{aligned}
A_{\perp L,R} &= \mathcal{N}_{K^*}(q^2) \sqrt{2} \lambda_{K^*}^{1/2} \left[[(C_9 - \mathcal{B}_{K^*} C_L^{(8)}) \mp (C_{10} - \mathcal{B}_{K^*} C_{L5}^{(8)})] \frac{V(q^2)}{m_B + m_{K^*}} + \frac{2m_b}{q^2} C_7 T_1(q^2) \right], \\
A_{\parallel L,R} &= -\mathcal{N}_{K^*}(q^2) \sqrt{2} (m_B^2 - m_{K^*}^2) \left[[(C_9 - \mathcal{B}_{K^*} C_L^{(8)}) \mp (C_{10} - \mathcal{B}_{K^*} C_{L5}^{(8)})] \frac{A_1(q^2)}{m_B - m_{K^*}} + \frac{2m_b}{q^2} C_7 T_2(q^2) \right], \\
A_{0L,R} &= -\frac{\mathcal{N}_{K^*}(q^2)}{2m_{K^*} \sqrt{q^2}} \left\{ [(C_9 - \mathcal{B}_{K^*} C_L^{(8)}) \mp (C_{10} - \mathcal{B}_{K^*} C_{L5}^{(8)})] \right. \\
&\quad \times \left[(m_B^2 - m_{K^*}^2 - q^2)(m_B + m_{K^*}) A_1(q^2) - \lambda_{K^*} \frac{A_2(q^2)}{m_B + m_{K^*}} \right] \\
&\quad \left. + 2m_b C_7 \left[(m_B^2 + 3m_{K^*}^2 - q^2) T_2(q^2) - \frac{\lambda_{K^*}}{m_B^2 - m_{K^*}^2} T_3(q^2) \right] \right\} + A_{0L,R}^{\text{NP}}, \\
A_t &= \frac{\mathcal{N}_{K^*}(q^2)}{\sqrt{q^2}} \left\{ \lambda_{K^*}^{1/2} [2(C_{10} - \mathcal{B}_{K^*} C_{L5}^{(8)})] A_0(q^2) - \beta_\ell(q^2) \cos \theta_\ell \frac{(m_B - m_{K^*})}{m_{K^*}} \right. \\
&\quad \left. \times [(m_B^2 - m_{K^*}^2 - q^2)(m_B + m_{K^*})^2 A_1(q^2) - \lambda_{K^*} A_2(q^2)] C_{LL}^{(8)} \right\}, \\
A_S &= -2\mathcal{N}_{K^*}(q^2) \lambda_{K^*}^{1/2} m_b m_\ell (C_S^{(8)} - C_{S'}^{(8)}) A_0(q^2),
\end{aligned} \tag{36}$$

where $\mathcal{B}_{K^*} = \beta_\ell \sqrt{\lambda_{K^*}} \cos \theta_\ell$, with $\lambda_{K^*} = m_B^4 + m_{K^*}^4 + q^4 - 2(m_B^2 m_{K^*}^2 + m_{K^*}^2 q^2 + m_B^2 q^2)$,

$$\mathcal{N}_{K^*}(q^2) = V_{tb} V_{ts}^* \left[\frac{G_F^2 \alpha^2}{3 \times 2^{10} \pi^5 m_B^3} q^2 \lambda_{K^*}^{1/2} \beta_\ell(q^2) \right]^{1/2},$$

and

$$\begin{aligned}
A_{0L}^{\text{NP}} &= \frac{\mathcal{N}_{K^*}(q^2)}{m_{K^*}} \sqrt{\frac{\lambda_{K^*}}{q^2} \beta_\ell(q^2) \cos \theta_\ell} \left[(m_B^2 - m_{K^*}^2 - q^2)(m_B + m_{K^*}) A_1(q^2) - \lambda_{K^*} \frac{A_2(q^2)}{m_B + m_{K^*}} \right] C_{LL}^{(8)}, \\
A_{0R}^{\text{NP}} &= 0.
\end{aligned} \tag{37}$$

We point out that Wilson coefficients $C_{L,L5}^{(8)}$ and $C_{LL}^{(8)}$ in our case are always accompanied by $\cos \theta_\ell$. In the limit of massless leptons and considering only the interference of $C_{L,L5}^{(8)}, C_{LL}^{(8)}$, with $C_{7,9,10}$ (i.e., neglecting higher powers of $C_i^{(8)}$ coefficients), all the angular functions I_i 's are of the form $I_i^{\text{SM}} + C_{7,9,10}(C_{L,L5}^{(8)}, C_{LL}^{(8)}) \cos \theta_\ell \approx I_i^{\text{SM}} + I_i^{\text{int}} \cos \theta_\ell$. The I_i^{SM} are the values of the angular coefficients in the SM, whereas I_i^{int} denotes the interference term of spin-2 dimension-8 operators with the standard dimension-6 ones. Such dependence on $\cos \theta_\ell$ of the NP terms leads us to consider the forward-backward asymmetry in θ_ℓ ,

$$\begin{aligned}
A_{\text{FB}}^{K^*} &= \frac{1}{\Gamma} \int dq^2 \int_0^{2\pi} d\phi \int_{-1}^1 d \cos \theta_{K^*} \left(\int_0^1 - \int_{-1}^0 \right) d \cos \theta_\ell \left(\frac{d^4 \Gamma}{dq^2 d \cos \theta_\ell d \cos \theta_{K^*} d\phi} \right) \\
&= \frac{3}{8} (2I_6^{s(\text{SM})} + I_6^{c(\text{int})} + I_1^{c(\text{int})} + 2I_1^{s(\text{int})}).
\end{aligned} \tag{38}$$

We present in Fig. 9 the branching ratio for $B \rightarrow K^* \mu^+ \mu^-$ in the low-energy region $1 \leq q^2 \leq 6 \text{ GeV}^2$, and we present in Fig. 10 the forward-backward asymmetry, as a function of the mass of spin-2 particle, respecting all relevant constraints from the $B_s - \bar{B}_s$ mixing, the top decay width, and muon anomalous magnetic moment. The spin-2 mediator increases the forward-backward asymmetry for the decay $B \rightarrow K^* \mu^+ \mu^-$ relative to the SM value $A_{\text{FB,SM}}^{K^*} = -0.035_{-0.033}^{+0.036}$ [74]. For the

FIG. 9. $\text{Br}(B \rightarrow K^* \mu^+ \mu^-)$ at $1 \leq q^2 \leq 6 \text{ GeV}^2$ as a function of m_G for $\Lambda = 1000 \text{ GeV}$ in the presence of the spin-2 contribution (red), taking into account the constraints from the muon anomalous magnetic moment, $B_s - \bar{B}_s$ mixing, and the top decay width. The SM value for the BR is shown by a blue line.

FIG. 10. A_{FB} as a function of the spin-2 mass, taking into account the constraints coming from the $B_s - \bar{B}_s$ mixing, the top decay width, and muon anomalous magnetic moment for the spin-2 contribution (green) for $\Lambda = 1000 \text{ GeV}$. The SM value for the A_{FB} is shown by a blue line.

rather low values of the spin-2 particle mass $m_G \simeq 20 - 40 \text{ GeV}$, the generated forward-backward asymmetry can be close to 15%. As in the case of $B \rightarrow K \mu^+ \mu^-$, due to the insignificant increase of the branching ratio for $B \rightarrow K^* \mu^+ \mu^-$, inclusion of the spin-2 mediator cannot explain the $R_{K^{(*)}}$ problem. The spin-2 particle also modifies the triple forward-backward asymmetries for K^* and the leptons and also the double asymmetries in ϕ and θ_K ,

$$\begin{aligned} A_2^{K^*} &= \frac{1}{\Gamma} \int dq^2 \left(\int_{-\pi/2}^{\pi/2} - \int_{\pi/2}^{3\pi/2} \right) d\phi \left(\int_0^1 - \int_{-1}^0 \right) \\ &\quad \times d \cos \theta_{K^*} \left(\int_0^1 - \int_{-1}^0 \right) d \cos \theta_\ell \Gamma_{\text{tot}} \\ &= \frac{1}{\pi} (2I_4^{(\text{SM})} + I_5^{(\text{int})}), \end{aligned} \quad (39)$$

$$\begin{aligned} A_3^{K^*} &= \frac{1}{\Gamma} \int dq^2 \left(\int_{-\pi/2}^{\pi/2} - \int_{\pi/2}^{3\pi/2} \right) d\phi \left(\int_0^1 - \int_{-1}^0 \right) \\ &\quad \times d \cos \theta_{K^*} \int_{-1}^1 d \cos \theta_\ell \Gamma_{\text{tot}} \\ &= \frac{3}{8} (2I_5^{(\text{SM})} + I_4^{(\text{int})}), \end{aligned} \quad (40)$$

where $\Gamma_{\text{tot}} = d^4\Gamma / (dq^2 d \cos \theta_\ell d \cos \theta_{K^*} d\phi)$. We show in Fig. 11 the asymmetries $A_2^{K^*}, A_3^{K^*}$ as a function of the spin-2 mass by taking into account all the experimental constraints considered in the text.

IV. CONCLUSIONS

We have considered the effects of the spin-2 mediator in $b \rightarrow s \mu^+ \mu^-$ transitions. We assumed that only flavor-nonuniversal couplings to left-handed b and s quarks exist and that in the leptonic sector the spin-2 particle couples to

FIG. 11. $A_2^{K^*}, A_3^{K^*}$ as a function of the spin-2 mass, taking into account the constraints coming from the $B_s - \bar{B}_s$ mixing, the top decay width, and muon anomalous magnetic moment for the spin-2 contribution (green) for $\Lambda = 1000 \text{ GeV}$. The SM value for the $A_{2,3}^{K^*}$ is shown by a blue line.

muons only. Naturally, the $B_s - \bar{B}_s$ mixing and the muon anomalous magnetic moment enable us to constrain these couplings. It is interesting that these constraints are very severe, not allowing any significant effect in the observables for the $B \rightarrow K\mu^+\mu^-$ and $B \rightarrow K^*\mu^+\mu^-$ decays. The model contains derivative couplings of the spin-2 particle with fermions. The dimension-8 effective Lagrangian induced by the spin-2 mediator creates additional terms in the angular distributions of the $B \rightarrow K^{(*)}\mu^+\mu^-$ decays. Such contributions are of the order $1/\Lambda^4$ and are therefore negligible. However, we find out that the spin-2 effective Lagrangian can create forward-backward asymmetries of the order $\sim 10\text{--}20\%$, for a rather small mass of the spin-2 particle $m_G \simeq 20\text{--}40$ GeV. This is especially interesting for $B \rightarrow K\mu^+\mu^-$ decay in which such an asymmetry vanishes in the SM.

Because of the left-handed nature of b , s , and μ in our model, one can also generate the coupling of the spin-2 particle to the top and charm quarks, $a_{bs}^L \simeq a_{tc}^L$. In that case, the spin-2 particle can be seen at the LHC, which would mean that for its mass higher than the top-quark mass the possible decay channels would be the top-quark and charm-quark jet, the b jet and light-quark jet, and the muon pair. A dedicated study of the spin-2 particle in the process $b\bar{s} \rightarrow \mu\mu$ as performed in Ref. [75] would put additional constraints on the couplings discussed in this paper.

The spin-2 mediation increases insignificantly the branching ratios $\text{BR}(B \rightarrow K^{(*)}\mu^+\mu^-)$, while the branching

ratios $\text{BR}(B \rightarrow K^{(*)}e^+e^-)$ are SM-like. This implies that our simple proposal cannot explain the observed R_K and R_{K^*} anomalies. One might expect that some full-fledged model containing additional scalars or gauge bosons together with a spin-2 particle can account for observed anomalies and also generate the forward-backward asymmetry.

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APPENDIX A: MATRIX ELEMENTS OF THE OPERATORS IN THE $B_s - \bar{B}_s$ MASS DIFFERENCE

The operators contributing to the mass difference in Eq. (24) can be determined using the results of QCD running given in Ref. [58],

$$\begin{aligned}
 (\Delta M_s)_{m_b^2 \bar{Q}_1} &= \frac{2}{3} m_{B_s} f_{B_s}^2 B_1(m_b) m_b^2(m_b) \left(\frac{(a_{sb}^L)^2}{4\Lambda^2 m_G^2} \right) [0.865 - 0.017\eta] \eta^{0.286}, \\
 (\Delta M_s)_{m_b^2 \bar{Q}_2} &= \frac{5}{12} \left(\frac{m_{B_s}}{m_b(m_b) + m_s(m_b)} \right)^2 m_{B_s} f_{B_s}^2 B_2(m_b) m_b^2(m_b) \left(\frac{(a_{sb}^L)^2}{16\Lambda^2 m_G^2} \frac{8}{3} \right) \\
 &\quad \{ [(1.879 - 0.18\eta)\eta^{-0.692} + (0.012 - 0.003\eta)\eta^{0.787}] \\
 &\quad - [(-0.493 - 0.014\eta)\eta^{-0.692} + (0.18 + 0.008\eta)\eta^{0.787}] \}, \\
 (\Delta M_s)_{m_b^2 \bar{Q}_3} &= -\frac{1}{12} \left(\frac{m_{B_s}}{m_b(m_b) + m_s(m_b)} \right)^2 m_{B_s} f_{B_s}^2 B_3(m_b) m_b^2(m_b) \left(\frac{(a_{sb}^L)^2}{16\Lambda^2 m_G^2} \frac{8}{3} \right) \\
 &\quad \{ [(-0.044 + 0.005\eta)\eta^{-0.692} + (0.035 - 0.012\eta)\eta^{0.787}] \\
 &\quad - [(0.011 + 0.0\eta)\eta^{-0.692} + (0.54 + 0.028\eta)\eta^{0.787}] \},
 \end{aligned} \tag{A1}$$

with $\eta = \alpha_S(\Lambda)/\alpha_S(m_t)$.

APPENDIX B: $B \rightarrow K\mu^+\mu^-$ AND $B \rightarrow K^*\mu^+\mu^-$ DECAYS

In the calculation, we use the following kinematics of $b \rightarrow s\mu^+\mu^-$ decay. The lepton pair is taken to be at rest, $\mathbf{k}_+ + \mathbf{k}_- = 0$, meaning that

$$k_+ + k_- = (q, 0, 0, 0), \quad k_- = \frac{q}{2}(1, \beta_\ell \sin \theta_\ell, 0, \beta_\ell \cos \theta_\ell), \quad k_+ = \frac{q}{2}(1, -\beta_\ell \sin \theta_\ell, 0, -\beta_\ell \cos \theta_\ell),$$

and

$$\begin{aligned}
p_b &= (E_b, 0, 0, \vec{p}_b), \\
p_s &= (E_s, 0, 0, \vec{p}_s), \\
\vec{p}_b &= \vec{p}_s = \frac{1}{2q} \sqrt{(m_b^2 - m_s^2 + q^2)^2 - 4q^2 m_b^2}, \\
E_b &= \frac{1}{2q} (m_b^2 + q^2 - m_s^2), \\
E_s &= \frac{1}{2q} (m_b^2 - m_s^2 - q^2), \\
\beta_\ell &= \sqrt{1 - 4 \frac{m_\ell^2}{q^2}},
\end{aligned} \tag{B1}$$

where θ_ℓ is the angle between the directions of b and ℓ^- .

The form factors for the $B \rightarrow K^*$ weak transitions are defined as

$$\begin{aligned}
\langle K^*(p_{K^*}) | \bar{s} \gamma_\mu P_{L,R} b | B(p) \rangle \\
= i \epsilon_{\mu\nu\alpha\beta} \epsilon^{\nu*} p^\alpha q^\beta \frac{V(q^2)}{m_B + m_{K^*}} \mp \frac{1}{2} \left\{ \epsilon_\mu^* (m_B + m_{K^*}) A_1(q^2) \right. \\
- (\epsilon^* \cdot q) (2p - q)_\mu \frac{A_2(q^2)}{m_B + m_{K^*}} \\
\left. - \frac{2m_{K^*}}{q^2} (\epsilon^* \cdot q) q_\mu [A_3(q^2) - A_0(q^2)] \right\},
\end{aligned} \tag{B2}$$

where

$$A_3(s) = \frac{m_B + m_{K^*}}{2m_{K^*}} A_1(q^2) - \frac{m_B - m_{K^*}}{2m_{K^*}} A_2(q^2), \tag{B3}$$

and

$$\begin{aligned}
\langle K^*(p_{K^*}) | \bar{s} i \sigma_{\mu\nu} q^\nu P_{R,L} b | B(p) \rangle \\
= -i \epsilon_{\mu\nu\alpha\beta} \epsilon^{\nu*} p^\alpha q^\beta T_1(q^2) \pm \frac{1}{2} \left\{ [\epsilon_\mu^* (m_B^2 - m_{K^*}^2) \right. \\
- (\epsilon^* \cdot q) (2p - q)_\mu] T_2(q^2) \\
\left. + (\epsilon^* \cdot q) \left[q_\mu - \frac{q^2}{m_B^2 - m_{K^*}^2} (2p - q)_\mu \right] T_3(q^2) \right\}.
\end{aligned} \tag{B4}$$

In the angular distribution given in Eq. (35), the following combinations of the amplitudes are introduced:

$$\begin{aligned}
I_1^s &= \frac{(2 + \beta_\ell^2)}{4} [|A_\perp^L|^2 + |A_\parallel^L|^2 + (L \rightarrow R)] \\
&\quad + \frac{4m_\ell^2}{q^2} \text{Re}(A_\perp^L A_\perp^{R*} + A_\parallel^L A_\parallel^{R*}), \\
I_1^c &= |A_0^L|^2 + |A_0^R|^2 + \frac{4m_\ell^2}{q^2} [|A_t|^2 + 2\text{Re}(A_0^L A_0^{R*})] + \beta_\ell^2 |A_S|^2, \\
I_2^s &= \frac{\beta_\ell^2}{4} [|A_\perp^L|^2 + |A_\parallel^L|^2 + (L \rightarrow R)], \\
I_2^c &= -\beta_\ell^2 [|A_0^L|^2 + (L \rightarrow R)], \\
I_3 &= \frac{1}{2} \beta_\ell^2 [|A_\perp^L|^2 - |A_\parallel^L|^2 + (L \rightarrow R)], \\
I_4 &= \frac{1}{\sqrt{2}} \beta_\ell^2 [\text{Re}(A_0^L A_\parallel^{L*}) + (L \rightarrow R)], \\
I_5 &= \sqrt{2} \beta_\ell \left[\text{Re}(A_0^L A_\perp^{L*}) - (L \rightarrow R) \right. \\
&\quad \left. - \frac{m_\ell}{\sqrt{q^2}} \text{Re}(A_\parallel^L A_S^* + A_\parallel^R A_S^*) \right], \\
I_6^s &= 2\beta_\ell [\text{Re}(A_\parallel^L A_\perp^{L*}) - (L \rightarrow R)], \\
I_6^c &= 4\beta_\ell \frac{m_\ell}{\sqrt{q^2}} \text{Re}[A_0^L A_S^* + (L \rightarrow R)], \\
I_7 &= \sqrt{2} \beta_\ell \left[\text{Im}(A_0^L A_\parallel^{L*}) - (L \rightarrow R) \right. \\
&\quad \left. + \frac{m_\ell}{\sqrt{q^2}} \text{Im}(A_\perp^L A_S^* + A_\perp^R A_S^*) \right], \\
I_8 &= \frac{1}{\sqrt{2}} \beta_\ell^2 [\text{Im}(A_0^L A_\perp^{L*}) + (L \rightarrow R)], \\
I_9 &= \beta_\ell^2 [\text{Im}(A_\parallel^{L*} A_\perp^L) + (L \rightarrow R)].
\end{aligned} \tag{B5}$$

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