Radiative and semileptonic B decays involving the tensor meson $K_2*(1430)$ in the standard model and beyond

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We study semileptonic and radiative B decays involving the strange tensor meson $K_2^*(1430)$ in the final state. Using the large energy effective theory (LEET) techniques, we formulate the $B \to K_2^*$ transition form factors in large recoil region. All the form factors can be parametrized in terms of two independent LEET functions $\zeta_\perp$ and $\zeta_\parallel$. The magnitude of $\zeta_\perp$ is estimated from the data for $\mathcal{B}(B \to K_2^*(1430)\gamma)$. Assuming a dipole q^2 dependence for the LEET functions and $\zeta_\parallel/\zeta_\perp = 1.0 \pm 0.2$, for which the former consists with the QCD counting rules, and the latter is favored by the $B \to \phi K_2^*$ data, we investigate the decays $B \to K_2^* \ell^+ \ell^-$ and $B \to K_2^* \nu\bar{\nu}$, where the contributions due to $\zeta_\parallel$ are suppressed by $m_{K_2^*}/m_B$. For the $B \to K_2^* \ell^+ \ell^-$ decay, in the large recoil region where the hadronic uncertainties are considerably reduced, the longitudinal distribution dF_L/ds is reduced by 20–30% due to the flipped sign of $c_7^{\rm eff}$ compared with the standard model result. Moreover, the forward-backward asymmetry zero is about 3.4 GeV^2 in the standard model, but changing the sign of $c_7^{\rm eff}$ yields a positive asymmetry for all values of the invariant mass of the lepton pair. We calculate the branching fraction for $B \to K_2^* \nu\bar{\nu}$ in the standard model. Our result exhibits the impressed resemblance between $B \to K_2^*(1430)\ell^+\ell^-$, $\nu\bar{\nu}$, and $B \to K^*(892)\ell^+\ell^-$, $\nu\bar{\nu}$.

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I. INTRODUCTION

The flavor-changing neutral current (FCNC) processes involving $b \to s(d)$ transitions occur only at loop level in the standard model (SM) and thus provide an important testing ground to look for new physics phenomena. Radiative B decays can offer bounds on the CKM matrix elements $|V_{ts}|$ and $|V_{td}|$ as well as powerful constraints on new physics. The absolute value of $c_7^{\rm eff}$, which is the Wilson coefficient of electromagnetic dipole operator, extracted from the current $B \to X_s\gamma$ data is consistent with the SM prediction within errors.

The $b \to s\ell^+\ell^-$ processes arise from photonic penguin, Z penguin, and W-box diagrams. The inclusive $B \to X_s\ell^+\ell^-$ and exclusive $B \to K^{(*)}\ell^+\ell^-$ decays have been measured [1,2]. We summarize the current data for branching fractions of exclusive radiative and semileptonic B decays relevant to the FCNC $b \to s$ transition in Table I [3–15]. The FCNC processes may receive sizable new-physics contributions [16–21]. Recently, *BABAR* and Belle measured interesting observables, K^* longitudinal fraction, forward-backward asymmetry, and isospin asymmetry, in the $B \to K^*\ell^+\ell^-$ decays [1,2,8,9,13,15]. Although the data are consistent with the SM predictions, all measurements favor the flipped-sign $c_7^{\rm eff}$ models [22]. The minimal flavor violation supersymmetry models with large $\tan\beta$ can be fine-tuned to have the flipped sign of $c_7^{\rm eff}$ [23,24], for which the charged Higgs is dominant. However, the contributions of the charged Higgs exchange to c_9 and c_{10} are suppressed by $1/\tan^2\beta$ for large $\tan\beta$.

The measurements of inclusive and various exclusive decays relevant to FCNC transitions can shed light on new physics. We have studied $B \to K_1(1270)\gamma$ and $B \to K_1(1270)\ell^+\ell^-$ [25,26], where the $K_1(1270)$ is the P-wave meson. $B \to K_1(1270)\gamma$ has been measured by Belle [27]. In this paper, we focus on the exclusive processes $B \to K_2^*(1430)\gamma$, $B \to K_2^*(1430)\ell^+\ell^-$, and $B \to K_2^*(1430)\nu\bar{\nu}$, where $K_2^*(1430)$ is the strange tensor meson with positive parity.

The $B \to K_2^*(1430)\gamma$ decays have been observed by the Belle and *BABAR* collaborations [6,7]. See also Table I. Corresponding semileptonic decays can be expected to be seen soon. Because both K_2^* and K^* mainly decay to the two-body $K\pi$ mode, therefore the angular-distribution analysis for the $B \to K^*\ell^+\ell^-$ decays are applicable to the study for $B \to K_2^*\ell^+\ell^-$ decays.

In experiments, the exclusive mode is much more easier to accessible than the inclusive process. However, the former contains form factors parametrizing hadronic matrix elements, and thus suffers from large theoretical uncertainties. $B \to K_2^*$ transition form factors, which are relevant to the study of the radiative and semileptonic B decays into a K_2^*, are less understood compared with $B \to K^*$ ones. So far there are only some quark model results about them [28–30]. In this paper we formulate the $B \to K_2^*$ form factors in the large recoil region using the large energy effective theory (LEET) techniques [31]. We will show that all the form factors can be parameterized in terms of two independent form factors $\zeta_\perp$ and $\zeta_\parallel$ in the LEET limit. The former form factor can be estimated by using the data for $B \to K_2^*(1430)\gamma$, while the latter only gives corrections of order $m_{K_2^*}/m_B$ in the amplitude.

TABLE I. Branching fractions of radiative and semileptonic B decays involving K^* or K_2^*.

Mode	$\mathcal{B}$ $[10^{-6}]$	Mode	$\mathcal{B}$ $[10^{-6}]$
$B^+ \to K^{*+}(892)\gamma$	45.7 ± 1.9 [3–5]	$B^0 \to K^{*0}(892)\gamma$	44.0 ± 1.5 [3–5]
$B^+ \to K_2^{*+}(1430)\gamma$	14.5 ± 4.3 [6]	$B^0 \to K_2^{*0}(1430)\gamma$	12.4 ± 2.4 [6,7]
$B^+ \to K^{*+}(892)e^+e^-$	$1.42^{+0.43}_{-0.39}$ [8,9]	$B^0 \to K^{*0}(892)e^+e^-$	$1.13^{+0.21}_{-0.18}$ [8,9]
$B^+ \to K^{*+}(892)\mu^+\mu^-$	$1.12^{+0.32}_{-0.27}$ [8–10]	$B^0 \to K^{*0}(892)\mu^+\mu^-$	$1.00^{+0.15}_{-0.13}$ [8–10]
$B^+ \to K^{*+}(892)\nu\bar{\nu}$	<80 [11,12]	$B^0 \to K^{*0}(892)\nu\bar{\nu}$	<120 [11,12]

We study the longitudinal distribution dF_L/ds and forward-backward asymmetry for the $B \to K_2^*\ell^+\ell^-$ decay. Particularly, we find that in the large recoil region, where the uncertainties of these observables arising from the form factors are considerably reduced not only due to taking the ratio of form factors but also due to the evaluation in the large $E_{K_2^*}$ limit. For the new-physics effect, we will focus on the possible correction due to the $c_7^{\rm eff}$ with the sign flipped.

We calculate the branching fraction for $B \to K_2^*\nu\bar{\nu}$ in the SM. This mode enhanced by the summation over three light neutrinos is theoretically cleaner due to the absence of long-distance corrections related to the relevant four-fermion operators. This decay is relevant for the nonstandard Z^0 coupling [32], light dark matter [33], and unparticles [34,35].

The paper is organized as follows: In Sec. II, we formulate the $B \to T$ form factors using the LEET techniques. In Sec. III, we numerically study the radiative and semileptonic B meson decays into the $K_2^*(1430)$. We conclude with a summary in Sec. IV.

II. $B \to T$ FORM FACTORS IN THE LEET

For simplicity we work in the rest frame of the B meson (with mass m_B) and assume that the light tensor meson T (with mass m_T) moves along the z axis. The momenta of the B and T are given by

$$p_B^\mu = (m_B, 0, 0, 0) \equiv m_B v^\mu, \qquad p_T^\mu = (E, 0, 0, p_3) \equiv E n^\mu, \tag{1}$$

respectively. Here, the tensor meson's energy E is given by

$$E = \frac{m_B}{2}\left(1 - \frac{q^2}{m_B^2} + \frac{m_T^2}{m_B^2}\right), \tag{2}$$

where $q \equiv p_B - p_T$. In the LEET limit,

$$E, m_B \gg m_T, \Lambda_{\rm QCD}, \tag{3}$$

we simply have

$$v^\mu = (1, 0, 0, 0), \qquad n^\mu \simeq (1, 0, 0, 1). \tag{4}$$

The polarization tensors $\varepsilon(\lambda)^{\mu\nu}$ of the massive spin-2 tensor meson with helicity λ can be constructed in terms of the polarization vectors of a massive vector state [36]

$$\varepsilon(0)^{*\mu} = (p_3, 0, 0, E)/m_T, \qquad \varepsilon(\pm)^{*\mu} = (0, \mp 1, +i, 0)/\sqrt{2}, \tag{5}$$

and are given by

$$\varepsilon^{\mu\nu}(\pm 2) \equiv \varepsilon(\pm)^\mu \varepsilon(\pm)^\nu, \tag{6}$$

$$\varepsilon^{\mu\nu}(\pm 1) \equiv \sqrt{\frac{1}{2}}(\varepsilon(\pm)^\mu \varepsilon(0)^\nu + \varepsilon(0)^\mu \varepsilon(\pm)^\nu), \tag{7}$$

$$\varepsilon^{\mu\nu}(0) \equiv \sqrt{\frac{1}{6}}(\varepsilon(+)^\mu \varepsilon(-)^\nu + \varepsilon(-)^\mu \varepsilon(+)^\nu) + \sqrt{\frac{2}{3}}\varepsilon(0)^\mu \varepsilon(0)^\nu. \tag{8}$$

Because of the purpose of the present study, we calculate the $\bar{B} \to T$ transition form factors

$$\langle T|V^\mu|\bar{B}\rangle, \qquad \langle T|A^\mu|\bar{B}\rangle, \qquad \langle T|T^{\mu\nu}|\bar{B}\rangle, \qquad \langle T|T_5^{\mu\nu}|\bar{B}\rangle, \tag{9}$$

where $V^\mu \equiv \bar{\psi}\gamma^\mu b$, $A^\mu \equiv \bar{\psi}\gamma^\mu\gamma_5 b$, $T^{\mu\nu} = \bar{\psi}\sigma^{\mu\nu}b$, and $T_5^{\mu\nu} = \bar{\psi}\sigma^{\mu\nu}\gamma_5 b$. There is a trick to write down the form factors in the LEET limit. We first note that we have three independent classes of Lorentz structures (i) $\epsilon^{\alpha\beta\mu\nu}$, (ii) v^μ, n^μ, and (iii)

$$\sqrt{2}\frac{m_T}{E}\{\varepsilon(\lambda)^{*\mu\nu}v_\nu - [\varepsilon(\lambda)^*_{\alpha\beta}v^\alpha v^\beta]n^\mu\} = \begin{cases} 0 & \text{for } \lambda = \pm 2, \\ \varepsilon(\pm)^\mu & \text{for } \lambda = \pm 1, \\ 0 & \text{for } \lambda = 0, \end{cases} \tag{10}$$

$$\sqrt{2}\frac{m_T}{E}\epsilon^{\mu\nu\rho\sigma}[\varepsilon(\lambda)^*_{\nu\alpha}v^\alpha]n_\rho v_\sigma = \begin{cases} 0 & \text{for } \lambda = \pm 2, \\ \epsilon^{\mu\nu\rho\sigma}\varepsilon(\pm)_\nu n_\rho v_\sigma & \text{for } \lambda = \pm 1, \\ 0 & \text{for } \lambda = 0, \end{cases} \tag{11}$$

$$\sqrt{\frac{3}{2}}\left(\frac{m_T}{E}\right)^2[\varepsilon(\lambda)^*_{\alpha\beta}v^\alpha v^\beta]n^\mu = \begin{cases} 0 & \text{for } \lambda = \pm 2, \\ 0 & \text{for } \lambda = \pm 1, \\ n^\mu & \text{for } \lambda = 0, \end{cases} \tag{12}$$

$$\sqrt{\frac{3}{2}}\left(\frac{m_T}{E}\right)^2[\varepsilon(\lambda)^*_{\alpha\beta}v^\alpha v^\beta]v^\mu = \begin{cases} 0 & \text{for } \lambda = \pm 2, \\ 0 & \text{for } \lambda = \pm 1, \\ v^\mu & \text{for } \lambda = \pm 0, \end{cases} \quad (13)$$

to project the relevant polarization states of the tensor meson. Equations (10), (12), and (13), are the vectors, but Eq. (11) is the axial vector. Matching the parities of the matrix elements and using the three classes of the Lorentz structures, we can then easily parametrize the form factors in the following results:

$$\langle T|V^\mu|\bar{B}\rangle = -i2E\left(\frac{m_T}{E}\right)\zeta_\perp^{(v)}\varepsilon^{*\mu\nu\rho\sigma}v_\nu n_\rho \varepsilon^*_{\sigma\beta}v^\beta, \quad (14)$$

$$\langle T|A^\mu|\bar{B}\rangle = 2E\left(\frac{m_T}{E}\right)\zeta_\perp^{(a)}[\varepsilon^{*\mu\alpha}v_\alpha - (\varepsilon^*_{\alpha\beta}v^\alpha v^\beta)n^\mu] + 2E\left(\frac{m_T^2}{E^2}\right)(\varepsilon^*_{\alpha\beta}v^\alpha v^\beta)[\zeta_\parallel^{(a)}n^\mu + \zeta_{\parallel,1}^{(a)}v^\mu], \quad (15)$$

$$\langle T|T^{\mu\nu}|\bar{B}\rangle = 2E\left(\frac{m_T^2}{E^2}\right)\zeta_\parallel^{(t)}\epsilon^{\mu\nu\rho\sigma}(\varepsilon^*_{\alpha\beta}v^\alpha v^\beta)v_\rho n_\sigma + 2E\left(\frac{m_T}{E}\right)\zeta_\perp^{(t)}\epsilon^{\mu\nu\rho\sigma}n_\rho[\varepsilon^*_{\sigma\alpha}v^\alpha - (\varepsilon^*_{\alpha\beta}v^\alpha v^\beta)n_\sigma] + 2E\left(\frac{m_T}{E}\right)\zeta_{\perp,1}^{(t)}\epsilon^{\mu\nu\rho\sigma}v_\rho \times[\varepsilon^*_{\sigma\alpha}v^\alpha - (\varepsilon^*_{\alpha\beta}v^\alpha v^\beta)n_\sigma], \quad (16)$$

$$\langle T|T_5^{\mu\nu}|\bar{B}\rangle = -i2E\left(\frac{m_T}{E}\right)\zeta_{\perp,1}^{(t_5)}\{[\varepsilon^{*\mu\alpha}v_\alpha - (\varepsilon^*_{\alpha\beta}v^\alpha v^\beta)n^\mu]v^\nu - (\mu\leftrightarrow\nu)\} - i2E\left(\frac{m_T}{E}\right)\zeta_\perp^{(t_5)}\{[\varepsilon^{*\mu\alpha}v_\alpha - (\varepsilon^*_{\alpha\beta}v^\alpha v^\beta)n^\mu]n^\nu - (\mu\leftrightarrow\nu)\} - i2E\left(\frac{m_T^2}{E^2}\right)\zeta_\parallel^{(t_5)}(\varepsilon^*_{\alpha\beta}v^\alpha v^\beta)(n^\mu v^\nu - n^\nu v^\mu), \quad (17)$$

where $\epsilon^{0123} = -1$ is adopted. $\langle T|T^{\mu\nu}|\bar{B}\rangle$ is related to $\langle T|T_5^{\mu\nu}|\bar{B}\rangle$ by using the relation: $\sigma^{\mu\nu}\gamma_5\epsilon_{\mu\nu\rho\sigma} = 2i\sigma_{\rho\sigma}$. Note that for the tensor meson only the states with helicities ± 1 and 0 contribute to the $\bar{B}\to T$ transition in the LEET limit. $\zeta_\perp$'s are relevant to T with helicity $= \pm 1$, and $\zeta_\parallel$'s to T with helicity $= 0$.

In order to reduce the number of the independent $\bar{B}\to T$ form factors, we consider the effective current operator $\bar{q}_n\Gamma b_v$ (with $\Gamma = 1$, γ_5, γ^μ, $\gamma^\mu\gamma_5$, $\sigma^{\mu\nu}$, $\sigma^{\mu\nu}\gamma_5$) in the LEET limit, instead of the original one $\bar{q}\Gamma b$ [31]. Here, b_v and q_n satisfy $\not{v}b_v = b_v$, $\not{n}q_n = 0$, and $(\not{n}\not{v}/2)q_n = q_n$. Employing the Dirac identities

$$\frac{\not{v}\not{n}}{2}\gamma^\mu = \frac{\not{v}\not{n}}{2}(n^\mu\not{v} - i\epsilon^{\mu\nu\rho\sigma}v_\nu n_\rho\gamma_\sigma\gamma_5), \quad (18)$$

$$\frac{\not{v}\not{n}}{2}\sigma^{\mu\nu} = \frac{\not{v}\not{n}}{2}[i(n^\mu v^\nu - n^\nu v^\mu) - i(n^\mu\gamma^\nu - n^\nu\gamma^\mu)\not{v} - \epsilon^{\mu\nu\rho\sigma}v_\nu n_\rho\gamma_\sigma\gamma_5], \quad (19)$$

where $\epsilon^{0123} = -1$ is adopted, one can obtain the following relations:

$$\bar{q}_n b_v = v_\mu\bar{q}_n\gamma^\mu b_v, \quad (20)$$

$$\bar{q}_n\gamma^\mu b_v = n^\mu\bar{q}_n b_v - i\epsilon^{\mu\nu\rho\sigma}v_\nu n_\rho\bar{q}_n\gamma_\sigma\gamma_5 b_v, \quad (21)$$

$$\bar{q}_n\gamma^\mu\gamma_5 b_v = -n^\mu\bar{q}_n\gamma_5 b_v - i\epsilon^{\mu\nu\rho\sigma}v_\nu n_\rho\bar{q}_n\gamma_\sigma b_v, \quad (22)$$

$$\bar{q}_n\sigma^{\mu\nu}b_v = i[n^\mu v^\nu\bar{q}_n b_v - n^\mu\bar{q}_n\gamma^\nu b_v - (\mu\leftrightarrow\nu)] - \epsilon^{\mu\nu\rho\sigma}v_\rho n_\sigma\bar{q}_n\gamma_5 b_v, \quad (23)$$

$$\bar{q}_n\sigma^{\mu\nu}\gamma_5 b_v = i[n^\mu v^\nu\bar{q}_n\gamma_5 b_v + n^\mu\bar{q}_n\gamma^\nu\gamma_5 b_v - (\mu\leftrightarrow\nu)] - \epsilon^{\mu\nu\rho\sigma}v_\rho n_\sigma\bar{q}_n b_v. \quad (24)$$

Substituting the above results into Eqs. (14)–(17), we have

$$\zeta_\perp^{(v)} = \zeta_\perp^{(a)} = \zeta_\perp^{(t)} = \zeta_\perp^{(t_5)} \equiv \zeta_\perp, \quad (25)$$

$$\zeta_\parallel^{(a)} = \zeta_\parallel^{(t)} = \zeta_\parallel^{(t_5)} \equiv \zeta_\parallel, \quad (26)$$

$$\zeta_{\parallel,1}^{(a)} = \zeta_{\perp,1}^{(t_5)} = \zeta_{\perp,1}^{(t)} = 0, \quad (27)$$

and thus find that there are only two independent components, $\zeta_\perp(q^2)$ and $\zeta_\parallel(q^2)$, for the $B\to T$ transition in the LEET limit. In the full theory, the $\bar{B}(p_B)\to\bar{K}_2^*(p_{K_2^*},\lambda)$ form factors are defined as follows:

$$\langle\bar{K}_2^*(p_{K_2^*},\lambda)|\bar{s}\gamma^\mu b|\bar{B}(p_B)\rangle = -i\frac{2}{m_B + m_{K_2^*}}\tilde{V}^{K_2^*}(q^2)\epsilon^{\mu\nu\rho\sigma}p_{B\nu}p_{K_2^*\rho}e^*_\sigma, \quad (28)$$

$$\langle\bar{K}_2^*(p_{K_2^*},\lambda)|\bar{s}\gamma^\mu\gamma_5 b|\bar{B}(p_B)\rangle = 2m_{K_2^*}\tilde{A}_0^{K_2^*}(q^2)\frac{e^*\cdot p_B}{q^2}q^\mu + (m_B + m_{K_2^*})\tilde{A}_1^{K_2^*}(q^2)\times\left[e^{*\mu} - \frac{e^*\cdot p_B}{q^2}q^\mu\right] - \tilde{A}_2^{K_2^*}(q^2)\frac{e^*\cdot p_B}{m_B + m_{K_2^*}}\times\left[p_B^\mu + p_{K_2^*}^\mu - \frac{m_B^2 - m_{K_2^*}^2}{q^2}q^\mu\right], \quad (29)$$

$$\langle\bar{K}_2^*(p_{K_2^*},\lambda)|\bar{s}\sigma^{\mu\nu}q_\nu b|\bar{B}(p_B)\rangle = 2\tilde{T}_1^{K_2^*}(q^2)\epsilon^{\mu\nu\rho\sigma}p_{B\nu}p_{K_2^*\rho}e^*_\sigma, \quad (30)$$

$$\langle \bar{K}_2^*(p_{K_2^*}, \lambda)|\bar{s}\sigma^{\mu\nu}\gamma_5 q_\nu b|\bar{B}(p_B)\rangle = -i\tilde{T}_2^{K_2^*}(q^2)[(m_B^2 - m_{K_2^*}^2)e^{*\mu} - (e^*\cdot p_B)(p_B^\mu + p_{K_2^*}^\mu)] - i\tilde{T}_3^{K_2^*}(q^2)(e^*\cdot p_B)\Big[q^\mu - \frac{q^2}{m_B^2 - m_{K_2^*}^2}(p_B^\mu + p_{K_2^*}^\mu)\Big], \quad (31)$$

where $e^\mu \equiv \varepsilon^{\mu\nu}(p_{K_2^*}, \lambda)p_{B,\nu}/m_B$ corresponding to $\lambda = 0, \pm 1$. We have $e^\mu = (|\vec{p}_{K_2^*}|/m_{K_2^*})\tilde{\varepsilon}^\mu$, where $\tilde{\varepsilon}(0) = \sqrt{2/3}\varepsilon(0)$ and $\tilde{\varepsilon}(\pm 1) = \sqrt{1/2}\varepsilon(\pm 1)$. We thus normalize these form factors and obtain relations as follows:

$$\tilde{A}_0^{K_2^*}(q^2)\frac{|\vec{p}_{K_2^*}|}{m_{K_2^*}} \equiv A_0^{K_2^*}(q^2) \simeq \Big(1 - \frac{m_{K_2^*}^2}{m_B E}\Big)\zeta_\parallel(q^2) + \frac{m_{K_2^*}}{m_B}\zeta_\perp(q^2), \quad (32)$$

$$\tilde{A}_1^{K_2^*}(q^2)\frac{|\vec{p}_{K_2^*}|}{m_{K_2^*}} \equiv A_1^{K_2^*}(q^2) \simeq \frac{2E}{m_B + m_{K_2^*}}\zeta_\perp(q^2), \quad (33)$$

$$\tilde{A}_2^{K_2^*}(q^2)\frac{|\vec{p}_{K_2^*}|}{m_{K_2^*}} \equiv A_2^{K_2^*}(q^2) \simeq \Big(1 + \frac{m_{K_2^*}}{m_B}\Big)\Big[\zeta_\perp(q^2) - \frac{m_{K_2^*}}{E}\zeta_\parallel(q^2)\Big], \quad (34)$$

$$\tilde{V}^{K_2^*}(q^2)\frac{|\vec{p}_{K_2^*}|}{m_{K_2^*}} \equiv V^{K_2^*}(q^2) \simeq \Big(1 + \frac{m_{K_2^*}}{m_B}\Big)\zeta_\perp(q^2), \quad (35)$$

$$\tilde{T}_1^{K_2^*}(q^2)\frac{|\vec{p}_{K_2^*}|}{m_{K_2^*}} \equiv T_1^{K_2^*}(q^2) \simeq \zeta_\perp(q^2), \quad (36)$$

$$\tilde{T}_2^{K_2^*}(q^2)\frac{|\vec{p}_{K_2^*}|}{m_{K_2^*}} \equiv T_2^{K_2^*}(q^2) \simeq \Big(1 - \frac{q^2}{m_B^2 - m_{K_2^*}^2}\Big)\zeta_\perp(q^2), \quad (37)$$

$$\tilde{T}_3^{K_2^*}(q^2)\frac{|\vec{p}_{K_2^*}|}{m_{K_2^*}} \equiv T_3^{K_2^*}(q^2) \simeq \zeta_\perp(q^2) - \Big(1 - \frac{m_{K_2^*}^2}{m_B^2}\Big)\frac{m_{K_2^*}}{E}\zeta_\parallel(q^2), \quad (38)$$

where have used $|\vec{p}_{K_2^*}|/E \simeq 1$. Our results are consistent with Ref. [30]. Defining

$$\tilde{\varepsilon}(0)^\mu = \alpha_L \varepsilon(0)^\mu, \qquad \tilde{\varepsilon}(\pm 1)^\mu = \beta_T \varepsilon(\pm 1)^\mu, \quad (39)$$

we can easily generalize the studies of $B \to K^*\gamma$, $B \to K^*\ell^+\ell^-$ and $B \to K^*\nu\bar{\nu}$ to $B \to K_2^*\gamma$, $B \to K_2^*\ell^+\ell^-$ and $B \to K_2^*\nu\bar{\nu}$ processes. For the K^* cases, we have $\alpha_L = \beta_T = 1$, whereas for the K_2^* cases, we instead use $\alpha_L = \sqrt{2/3}$ and $\beta_T = 1/\sqrt{2}$.

III. NUMERICAL STUDY

In the following numerical study, we use the input parameters listed in Table II. The Wilson coefficients that we adopt are the same as that in Ref. [26].

A. $B \to K_2^*\gamma$ and $B \to K_2^*\ell^+\ell^-$

The effective Hamiltonian relevant to the $B \to K_2^*\gamma$ and $B \to K_2^*\ell^+\ell^-$ decays is given by

$$\mathcal{H}_{\rm eff} = -\frac{G_F}{\sqrt{2}}V_{tb}V_{ts}^*\sum_{i=1}^{10} c_i(\mu)\mathcal{O}_i(\mu) + {\rm H.c.}, \quad (40)$$

$$\begin{aligned} \mathcal{O}_7 &= -\frac{g_{\rm em}m_b}{8\pi^2}\bar{s}\sigma_{\mu\nu}(1+\gamma_5)bF_{\mu\nu},\\ \mathcal{O}_8 &= -\frac{g_s m_b}{8\pi^2}\bar{s}_i\sigma_{\mu\nu}(1+\gamma_5)b_j G^{\mu\nu}T^{ij},\\ \mathcal{O}_9 &= \frac{\alpha_{\rm EM}}{2\pi}\bar{s}(1-\gamma_5)b(\bar{\ell}\ell),\\ \mathcal{O}_{10} &= \frac{\alpha_{\rm EM}}{2\pi}\bar{s}(1-\gamma_5)b(\bar{\ell}\gamma_5\ell). \end{aligned} \quad (41)$$

In analogy to $B \to K^*\gamma$ [24,39–41], the $B \to K_2^*\gamma$ decay width reads

$$\Gamma(B \to K_2^*\gamma) = \frac{G_F^2\alpha_{\rm EM}|V_{ts}^*V_{tb}|^2}{32\pi^4}m_{b,\rm pole}^2 m_B^3\Big(1 - \frac{m_{K_2^*}^2}{m_B^2}\Big)^3 \times |c_7^{(0)\rm eff} + A^{(1)}|^2|T_1^{K_2^*}(0)|^2\beta_T^2, \quad (42)$$

with $\beta_T = \sqrt{1/2}$. Here, $A^{(1)}$ is decomposed into the following components [40]:

$$A^{(1)}(\mu) = A_{c_7}^{(1)}(\mu) + A_{\rm ver}^{(1)}(\mu) = -0.038 - 0.016i. \quad (43)$$

In the LEET limit, $T_1^{K_2^*}(q^2)$ can be parametrized in terms of two independent functions $\zeta_\perp(q^2)$ and $\zeta_\parallel(q^2)$. Using $c_7^{(0)\rm eff} = -0.315$ and the $\mathcal{B}(B^0 \to K_2^{*0}\gamma)$ data in Table I, we estimate the value of $\zeta_\perp(0)$ as

TABLE II. Input parameters

Tensor meson mass	$m_{K_2^{*+}(1430)} = 1.426$ GeV, $m_{K_2^{*0}(1430)} = 1.432$ GeV,
b quark mass [37]	$m_{b,\rm pole} = 4.79^{+0.19}_{-0.08}$ GeV,
B lifetime (picosecond)	$\tau_{B^+} = 1.638$, $\tau_{B^0} = 1.530$,
CKM parameter [38]	$|V_{ts}^*V_{tb}| = 0.040 \pm 0.001$

$$T_1^{K_2^*}(0) \simeq \zeta_\perp(0) = 0.27 \pm 0.03^{+0.00}_{-0.01}, \tag{44}$$

where the errors are due to the uncertainties of the experimental data and pole mass of the b quark, respectively. The uncertainty is mainly due to the error of the data. We use the QCD counting rules to analyze the q^2 dependence of form factors [42]. We consider the Breit frame, where the initial B meson moves in the opposite direction but with the same magnitude of the momentum compared with the final state K_2^*, i.e., $\vec{p}_B = -\vec{p}_{K_2^*}$. In the large recoil region, where $q^2 \sim 0$, since the two quarks in mesons have to interact strongly with each other to turn around the spectator quark, the transition amplitude is dominated by the one-gluon exchange between the quark pair and is therefore proportional to $1/E^2$. Thus, we get $\langle K_2^*(p_{K_2^*}, \pm 1)|V^\mu|B(p_B)\rangle \propto \epsilon^{\mu\nu\rho\sigma} p_{B\nu} p_{K_2^*\rho} \varepsilon(\pm)_\sigma \times 1/E^2$ and $\langle K_2^*(p_{K_2^*}, 0)|A^\mu|B(p_B)\rangle \propto p^\mu_{K_2^*} \times 1/E^2$. In other words, we have $\zeta_{\perp,\parallel}(q^2) \sim 1/E^2$ in the large recoil region. Motivated by the above analysis, we will model the q^2 dependence of the form-factor functions to be $\zeta_{\perp,\parallel}(q^2) = \zeta_{\perp,\parallel}(0) \cdot (1 - q^2/m_B^2)^{-2}$. For the value[1] of $\zeta_\parallel(0)$, within the framework of the SM model, it was shown that $f_T/f_L \approx 3(m_\phi/m_B)^2(\zeta_\perp/\zeta_\parallel)^2$ for the $B \to \phi K_2^*$ decay [30], where f_T and f_L are the transverse and longitudinal components, respectively.[2] Comparing with the current data $f_L = 0.80 \pm 0.10$ for $B^+ \to \phi K_2^*(1430)^+$ and $f_L = 0.901^{+0.059}_{-0.061}$ for $B^0 \to \phi K_2^*(1430)^0$ [43], we therefore parametrize

$$\xi \equiv \zeta_\parallel(0)/\zeta_\perp(0), \quad \text{with} \quad 0.8 \le \xi \le 1.2 \tag{45}$$

to take into account the possible uncertainty.

The invariant amplitude of $\bar{B} \to \bar{K}_2^* \ell^+ \ell^-$, in analogy to [24], is given by

$$\mathcal{M} = -i\frac{G_F \alpha_{\rm EM}}{2\sqrt{2}\pi} V_{ts}^* V_{tb} m_B [\mathcal{T}_\mu \bar{s}\gamma^\mu b + \mathcal{U}_\mu \bar{s}\gamma^\mu\gamma_5 b], \tag{46}$$

where

$$\mathcal{T}_\mu = \mathcal{A}\epsilon_{\mu\nu\rho\sigma}\tilde{\varepsilon}^{*\nu} p_B^\rho p_T^\sigma - i m_B^2 \mathcal{B}\tilde{\varepsilon}^*_\mu + i\mathcal{C}(\tilde{\varepsilon}^* \cdot p_B) p_\mu + i\mathcal{D}(\tilde{\varepsilon}^* \cdot p_B) q_\mu, \tag{47}$$

$$\mathcal{U}_\mu = \mathcal{E}\epsilon_{\mu\nu\rho\sigma}\tilde{\varepsilon}^{*\nu} p_B^\rho p_T^\sigma - i m_B^2 \mathcal{F}\tilde{\varepsilon}^*_\mu + i\mathcal{G}(\tilde{\varepsilon}^* \cdot p_B) p_\mu + i\mathcal{H}(\tilde{\varepsilon}^* \cdot p_B) q_\mu. \tag{48}$$

The $\mathcal{D}$ term vanishes when equations of motion of leptons are taken into account. The building blocks $\mathcal{A}, \cdots, \mathcal{H}$ are given by

$$\mathcal{A} = \frac{2}{1 + \hat{m}_{K_2^*}} c_9^{\rm eff} V^{K_2^*}(s) + \frac{4\hat{m}_b}{\hat{s}} c_7^{\rm eff} T_1^{K_2^*}(s), \tag{49}$$

$$\mathcal{B} = (1 + \hat{m}_{K_2^*})\Big[c_9^{\rm eff}(\hat{s}) A_1^{K_2^*}(s) + 2\frac{\hat{m}_b}{\hat{s}}(1 - \hat{m}_{K_2^*}) c_7^{\rm eff} T_2^{K_2^*}(s)\Big], \tag{50}$$

$$\mathcal{C} = \frac{1}{1 - \hat{m}_{K_2^*}}\Big[(1 - \hat{m}_{K_2^*}) c_9^{\rm eff}(\hat{s}) A_2^{K_2^*}(s) + 2\hat{m}_b c_7^{\rm eff}\Big(T_3^{K_2^*}(s) + \frac{1 - \hat{m}_{K_2^*}}{\hat{s}} T_2^{K_2^*}(s)\Big)\Big], \tag{51}$$

$$\mathcal{D} = \frac{1}{\hat{s}}[c_9^{\rm eff}(\hat{s})\{(1 + \hat{m}_{K_2^*}) A_1^{K_2^*}(s) - (1 - \hat{m}_{K_2^*}) A_2^{K_2^*}(s)\} - 2\hat{m}_{K_2^*} A_0^{K_2^*}(s) - 2\hat{m}_b c_7^{\rm eff} T_3^{K_2^*}(s)], \tag{52}$$

$$\mathcal{E} = \frac{2}{1 + \hat{m}_{K_2^*}} c_{10} V^{K_2^*}(s), \qquad \mathcal{F} = (1 + \hat{m}_{K_2^*}) c_{10} A_1^{K_2^*}(s), \qquad \mathcal{G} = \frac{1}{1 + \hat{m}_{K_2^*}} c_{10} A_2^{K_2^*}(s), \tag{53}$$

$$\mathcal{H} = \frac{1}{\hat{s}} c_{10}[(1 + \hat{m}_{K_2^*}) A_1^{K_2^*}(s) - (1 - \hat{m}_{K_2^*}) A_2^{K_2^*}(s) - 2\hat{m}_{K_2^*} A_0^{K_2^*}(s)], \tag{54}$$

where $\hat{s} \equiv s/m_B^2$ and $s \equiv (p_+ + p_-)^2$ with $p_\pm$ being the momenta of the leptons $\ell^\pm$. $c_9^{\rm eff}(\hat{s}) = c_9 + Y_{\rm pert}(\hat{s}) + Y_{\rm LD}$ contains both the perturbative part $Y_{\rm pert}(\hat{s})$ and long-distance part $Y_{\rm LD}(\hat{s})$. $Y(\hat{s})_{\rm LD}$ involves $B \to K_1 V(\bar{c}c)$ resonances, where $V(\bar{c}c)$ are the vector charmonium states. We follow Refs. [44,45] and set

$$Y_{\rm LD}(\hat{s}) = -\frac{3\pi}{\alpha_{\rm EM}^2} c_0 \sum_{V=\psi(1s),\cdots} \kappa_V \frac{\hat{m}_V \mathcal{B}(V \to \ell^+\ell^-)\hat{\Gamma}^V_{\rm tot}}{\hat{s} - \hat{m}_V^2 + i\hat{m}_V \hat{\Gamma}^V_{\rm tot}}, \tag{55}$$

where $\hat{\Gamma}^V_{\rm tot} \equiv \Gamma^V_{\rm tot}/m_B$ and $\kappa_V = 2.3$. The detailed parameters used in this paper can be found in Ref. [26]. The longitudinal, transverse, and total differential decay rates are, respectively, given by

$$\frac{d\Gamma_L}{ds} \equiv \frac{d\Gamma}{ds}\bigg|_{\substack{\alpha_L=\sqrt{2/3}\\ \beta_T=0}}, \qquad \frac{d\Gamma_T}{ds} \equiv \frac{d\Gamma}{ds}\bigg|_{\substack{\alpha_L=0\\ \beta_T=\sqrt{1/2}}}, \qquad \frac{d\Gamma_{\rm total}}{ds} \equiv \frac{d\Gamma}{ds}\bigg|_{\substack{\alpha_L=\sqrt{2/3}\\ \beta_T=\sqrt{1/2}}}, \tag{56}$$

with

[1]The light-front results infer that $\zeta_\perp$ and $\zeta_\parallel$ are of the same sign [28].

[2]Here, the new physics contribution can be negligible if it mainly affects $c_7^{\rm eff}$.

$$\frac{d\Gamma}{d\hat{s}} = \frac{G_F^2 \alpha_{\rm EM}^2 m_B^5}{2^{10}\pi^5}|V_{ts}^* V_{tb}|^2 \bigg\{ \frac{1}{6}|\mathcal{A}|^2 \hat{u}(\hat{s})\hat{s}\beta_T^2\{3[1-2(\hat{m}_{K_2^*}^2+\hat{s})+(\hat{m}_{K_2^*}^2-\hat{s})^2]-\hat{u}(\hat{s})^2\} + \beta_T^2|\mathcal{E}|^2\hat{s}\frac{\hat{u}(\hat{s})^3}{3}$$
$$+\frac{1}{12\hat{m}_{K_2^*}^2\lambda}|\mathcal{B}|^2\hat{u}(\hat{s})\{3[1-2(\hat{m}_{K_2^*}^2+\hat{s})+(\hat{m}_{K_2^*}^2-s)^2]-\hat{u}(\hat{s})^2\}[(-1+\hat{m}_{K_2^*}^2+\hat{s})^2\alpha_L^2+8\hat{m}_{K_2^*}^2\hat{s}\beta_T^2]$$
$$+\frac{1}{12m_{K_2^*}^2\lambda}|\mathcal{F}|^2\hat{u}(\hat{s})\{3\alpha_L^2\lambda^2+\hat{u}(\hat{s})^2[16\hat{m}_{K_2^*}^2\hat{s}\beta_T^2-(1-2(\hat{m}_{K_2^*}^2+\hat{s})+\hat{m}_{K_2^*}^4+\hat{s}^2-10\hat{m}_{K_2^*}^2\hat{s})\alpha_L^2]\}$$
$$+\alpha_L^2\hat{u}(s)\frac{\lambda}{4\hat{m}_{K_2^*}^2}\bigg[|\mathcal{C}|^2\bigg(\lambda-\frac{\hat{u}(\hat{s})^2}{3}\bigg)+|\mathcal{G}|^2\bigg(\lambda-\frac{\hat{u}(\hat{s})^2}{3}+4\hat{m}_\ell^2(2+2\hat{m}_{K_2^*}^2-\hat{s})\bigg)\bigg]$$
$$-\alpha_L^2\hat{u}(s)\frac{1}{2\hat{m}_{K_2^*}^2}\bigg[\mathrm{Re}(\mathcal{B}\mathcal{C}^*)\bigg(\lambda-\frac{\hat{u}(\hat{s})^2}{3}\bigg)(1-\hat{m}_{K_2^*}^2-\hat{s})+\mathrm{Re}(\mathcal{F}\mathcal{G}^*)\bigg\{\bigg(\lambda-\frac{\hat{u}(\hat{s})^2}{3}\bigg)(1-\hat{m}_{K_2^*}^2-\hat{s})+4\hat{m}_\ell^2\lambda\bigg\}\bigg]$$
$$-2\alpha_L^2\hat{u}(s)\frac{\hat{m}_\ell^2}{\hat{m}_{K_2^*}^2}\lambda[\mathrm{Re}(\mathcal{F}\mathcal{H}^*)-\mathrm{Re}(\mathcal{G}\mathcal{H}^*)(1-\hat{m}_{K_2^*}^2)]+\alpha_L^2\hat{u}(s)\frac{\hat{m}_\ell^2}{\hat{m}_{K_2^*}^2}\hat{s}\lambda|\mathcal{H}|^2\bigg\}. \quad (57)$$

We have chosen the kinematic variables $\hat{u} \equiv u/m_B^2$ and $\hat{u}s \equiv u(s)/m_B^2$, where $u = -u(s)\cos\theta$ and

$$u(s) \equiv \sqrt{\lambda\bigg(1-\frac{4\hat{m}_\ell^2}{\hat{s}}\bigg)}, \quad (58)$$

with

$$\lambda \equiv 1+\hat{m}_{K_2^*}^4+\hat{s}^2-2\hat{m}_{K_2^*}^2-2\hat{s}-2\hat{m}_{K_2^*}^2\hat{s}, \quad (59)$$

and θ being the angle between the moving direction of ℓ^+ and B meson in the center of mass frame of the $\ell^+\ell^-$ pair. In Fig. 1, the total decay rates for $B \to K_2^*(1430)\mu^+\mu^-$ with and without charmonium resonances are plotted. The detailed results for the charmonium resonances can be found in Refs. [44,45]. The branching fraction for non-resonant $B \to K_2^*\mu^+\mu^-$ is obtained to be

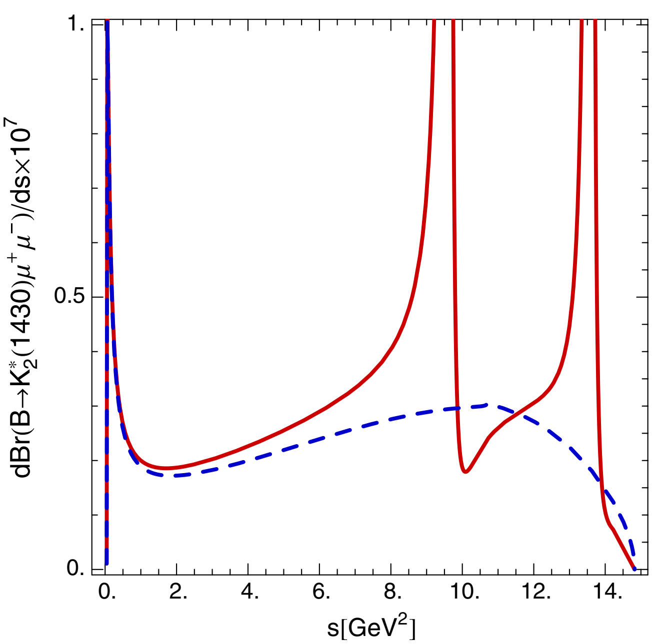

FIG. 1 (color online). The differential decay rates $d\Gamma_{\rm total}(B^0 \to K_2^{*0}(1430)\mu^+\mu^-)/ds$ as functions of the dimuon invariant mass s. The solid (dashed) curve corresponds to the center value of the decay rate with (without) the charmonium resonance effects.

$$\mathcal{B}(B^0 \to K_2^{*0}(1430)\mu^+\mu^-) = (3.5^{+1.1+0.7}_{-1.0-0.6}) \times 10^{-7}, \quad (60)$$

where the first error comes from the variation of $\zeta_\perp$ in Eq. (44), the second error from the uncertainty of ξ in Eq. (45).

The longitudinal fraction distribution for $B \to K_2^*\ell^+\ell^-$ decay is defined as

$$\frac{dF_L}{ds} \equiv \frac{d\Gamma_L}{ds}\bigg/\frac{d\Gamma_{\rm total}}{ds}. \quad (61)$$

In Fig. 2, the longitudinal fraction distribution for the $B \to K_2^*(1430)\mu^+\mu^-$ decay is plotted. For comparison, we also plot $F_L(B \to K^*(892)\mu^+\mu^-)/ds$ as a benchmark. For small s ($\lesssim 3$ GeV^2), $B \to K^*\mu^+\mu^-$, and $B \to K_2^*\mu^+\mu^-$ have similar rates for the longitudinal fraction, while for large s ($\gtrsim 4$ GeV^2) the dF_L/ds for the $B \to K_2^*\mu^+\mu^-$ decay slightly exceeds the $B \to K^*\mu^+\mu^-$. More interestingly, when $s \sim 3$ GeV^2, the result of the new physics models with the flipped-sign solution for $c_7^{\rm eff}$ can deviate more remarkably from the SM prediction (and can be reduced by 20–30%).

The forward-backward asymmetry for the $\bar{B} \to \bar{K}_2^*\ell^+\ell^-$ decay is given by

$$\frac{dA_{\rm FB}}{d\hat{s}} = -\beta_T^2\frac{G_F^2\alpha_{\rm EM}^2 m_B^5}{2^{10}\pi^5}|V_{ts}^*V_{tb}|^2\hat{s}\,\hat{u}(s)^2[\mathrm{Re}(\mathcal{B}\mathcal{E}^*) + \mathrm{Re}(\mathcal{A}\mathcal{F}^*)]\big|_{\beta_T=\sqrt{1/2}} \quad (62)$$

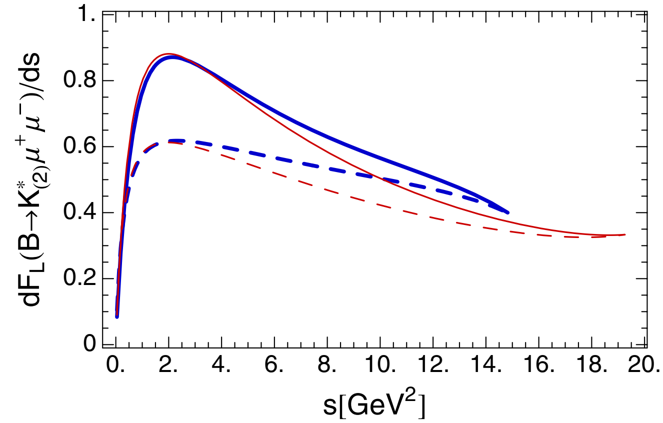

FIG. 2 (color online). Longitudinal fraction distributions dF_L/ds as functions of s. The thick (blue) and thin (red) curves correspond to the central values of $B \to K_2^{*0}(1430)\mu^+\mu^-$ and $B^0 \to K^{*0}(892)\mu^+\mu^-$ decays, respectively. The solid and dashed curves correspond to the SM and new physics model with the flipped sign of $c_7^{\rm eff}$, respectively.

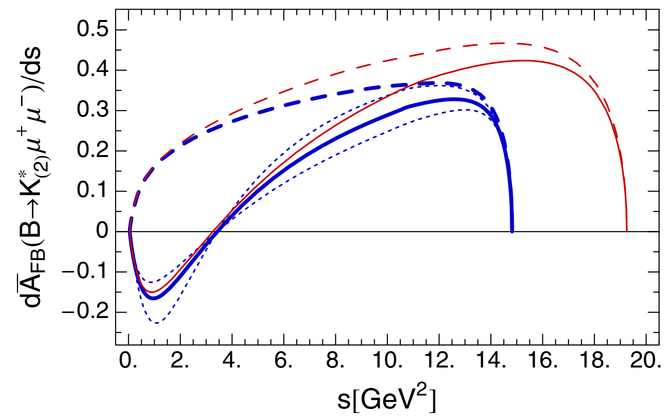

FIG. 3 (color online). Forward-backward asymmetries $d\bar{A}_{\rm FB}/ds$ for $B \to K_2^*(1430)\mu^+\mu^-$(thick curves) and $B \to K^*(892)\mu^+\mu^-$ (thin curves) as functions of the dimuon invariant mass s. The solid and dashed curves correspond to the SM and new physics model with the flipped sign of $c_7^{\rm eff}$. Variation due to the uncertainty from $\zeta_\parallel(q^2)/\zeta_\perp(q^2)$ [see Eq. (45)] is denoted by dotted curves.

$$= -\beta_T^2 \frac{G_F^2 \alpha_{\rm EM}^2 m_B^5}{2^{10}\pi^5} |V_{ts}^* V_{tb}|^2 \hat{s}\, \hat{u}(s)^2 \bigg[{\rm Re}(c_{10} c_9^{\rm eff}) V^{K_2^*} A_1^{K_2^*} + \frac{\hat{m}_b}{\hat{s}} {\rm Re}(c_{10} c_7^{\rm eff}) \{(1 - \hat{m}_{K_2^*}) V^{K_2^*} T_2^{K_2^*} + (1 + \hat{m}_{K_2^*}) A_1^{K_2^*} T_1^{K_2^*}\} \bigg] \Big|_{\beta_T = \sqrt{1/2}}. \tag{63}$$

In Fig. 3 we illustrate the normalized forward-backward asymmetry $d\bar{A}_{\rm FB}/ds \equiv (dA_{\rm FB}/ds)/(d\Gamma_{\rm total}/ds)$ for $B \to K_2^*\mu^+\mu^-$ together with $B \to K^*\mu^+\mu^-$.

In the SM, the forward-backward asymmetry zero s_0 for $B \to K_2^*\mu^+\mu^-$ is defined by

$$\begin{aligned} &{\rm Re}[c_{10} c_9^{\rm eff}(\hat{s}_0)] V^{K_2^*}(s_0) A_1^{K_2^*}(s_0) \\ &= -\frac{\hat{m}_b}{\hat{s}_0} {\rm Re}(c_{10} c_7^{\rm eff}) \{(1 - \hat{m}_{K_2^*}) V^{K_2^*}(s_0) T_2^{K_2^*}(s_0) \\ &\quad + (1 + \hat{m}_{K_2^*}) A_1^{K_2^*}(s_0) T_1^{K_2^*}(s_0)\}. \end{aligned} \tag{64}$$

We obtain

$$s_0 = 3.4 \pm 0.1 \ {\rm GeV}^2, \tag{65}$$

where the error comes from the variation of m_b. This result is very close to the zero for $B \to K^*\mu^+\mu^-$. As shown in Fig. 3, it is interesting to note that the form factor uncertainty of the zero vanishes in the LEET limit.

The asymmetry zero exists only for ${\rm Re}[c_9^{\rm eff}(s) c_{10}]{\rm Re}(c_7^{\rm eff} c_{10}) < 0$. Therefore, with the flipped sign of $c_7^{\rm eff}$ along, compared with the SM prediction, the asymmetry zero disappears, and $dA_{\rm FB}/ds$ is positive for all values of s. From recent measurements for $B \to K^*\ell^+\ell^-$ decays, the solution with the flipped sign of $c_7^{\rm eff}$ seems to be favored by the data [22,46,47]. One can find the further discussion in Ref. [26] for the $B \to K_1(1270)\ell^+\ell^-$ decays.

B. $B \to K_2^* \nu\bar{\nu}$

In the SM, $b \to s\nu\bar{\nu}$ proceeds through Z penguin and box diagrams involving top quark exchange [48]. One of the reasons that we are interested in the study of decays going through $b \to s\nu\bar{\nu}$ is the absence of long-distance corrections related to the relevant four-fermion operators. Moreover, the branching fractions are enhanced by the summation over three light neutrinos. New physics contributions arising from new loop and/or box diagrams may significantly modify the predictions. In the SM, the branching fractions involving K or K^* are predicted to be $\mathcal{B}(B \to K\nu\bar{\nu}) \simeq 3.8 \times 10^{-6}$ and $\mathcal{B}(B \to K^*\nu\bar{\nu}) \simeq 13 \times 10^{-6}$ [48,49], while only upper limits $10^{-4} \sim 10^{-5}$ were set in the experiments [11,12,22]. In the new physics scenario, the contribution originating from the nonstandard Z^0 coupling can enhance the branching fraction by a factor of 10 [32]. This mode is also relevant to search for light dark matter [33] and unparticles [34,35].

The generally effective weak Hamiltonian relevant to the $b \to s\nu\bar{\nu}$ decay is given by

$$\begin{aligned} \mathcal{H}_{\rm eff} &= c_L \bar{s}\gamma^\mu(1 - \gamma_5) b \bar{\nu}\gamma_\mu(1 - \gamma_5)\nu \\ &\quad + c_R \bar{s}\gamma^\mu(1 + \gamma_5) b \bar{\nu}\gamma_\mu(1 - \gamma_5)\nu, \end{aligned} \tag{66}$$

where c_L and c_R are left- and right-handed weak hadronic current contributions, respectively. New physics effects can modify the SM value of c_L, while c_R only receives the contribution from physics beyond the SM [32]. In the SM we have

$$c_L^{\rm SM} = \frac{G_F}{\sqrt{2}} \frac{\alpha_{\rm EM}}{2\pi \sin^2\theta_W} V_{tb} V_{ts}^* X(x_t) = 2.9 \times 10^{-9}, \tag{67}$$

where the detailed form of $X(x_t)$ can be found in Refs. [50–53]. The K_2^* helicity polarization rates of the missing

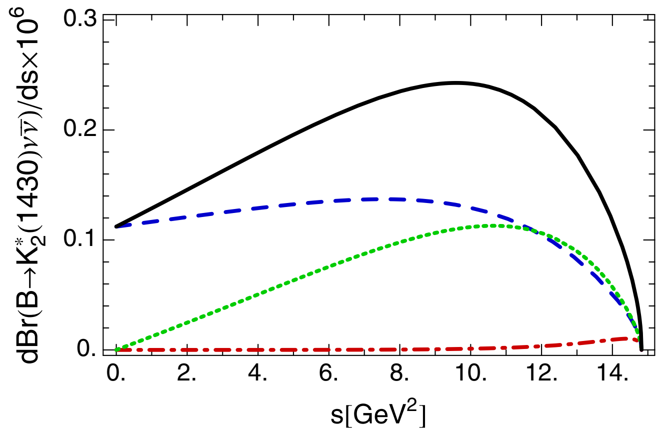

FIG. 4 (color online). Branching fraction distribution $d\mathcal{B}(B^0 \to K_2^{*0}\bar{\nu}\nu)/ds$ as a function of the missing invariant mass squared s within the SM. The solid (black), dashed (blue), dotted (green), and dotted-dashed (red) curves correspond to the total decay rate and the polarization rates with helicities $h = 0, -1, +1$, respectively.

invariant mass-squared distribution $d\Gamma_h/ds$ of the $B \to K_2^*\bar{\nu}\nu$ decay are given by [32,54–56],

$$\frac{d\Gamma_0}{d\hat{s}} = 3\alpha_L^2 \frac{|\vec{p}|}{48\pi^3} \frac{|c_L - c_R|^2}{m_{K_2^*}^2}\Big[(m_B + m_{K_2^*}) \times (m_B E - m_{K_2^*}^2)A_1^{K_2^*}(q^2) - \frac{2m_B^2}{m_B + m_{K_2^*}}|\vec{p}|^2 A_2^{K_2^*}(q^2)\Big]^2, \quad (68)$$

$$\frac{d\Gamma_{\pm 1}}{d\hat{s}} = 3\beta_T^2 \frac{|\vec{p}|q^2}{48\pi^3}\Big|(c_L + c_R)\frac{2m_B|\vec{p}|}{m_B + m_{K_2^*}}V^{K_2^*}(q^2) \mp (c_L - c_R)(m_B + m_{K_2^*})A_1^{K_2^*}(q^2)\Big|^2, \quad (69)$$

where $\hat{s} \equiv s/m_B^2$, $\alpha_L = \sqrt{2/3}$ and $\beta_T = \sqrt{1/2}$ with $0 \leq s \leq (m_B - m_{K_2^*})^2$ being the invariant mass squared of the neutrino-antineutrino pair. Here, the factor 3 counts the numbers of the neutrino generations. $\vec{p}$ and E are the three-momentum and energy of the K_2^* in the B rest frame. In Fig. 4, we show the distribution of the missing invariant mass squared for the $B \to K_2^*(1430)\bar{\nu}\nu$ decay within the SM. We find

$$\mathcal{B}(B^0 \to K_2^{*0}(1430)\bar{\nu}\nu) = (2.8^{+0.9+0.6}_{-0.8-0.5}) \times 10^{-6}, \quad (70)$$

where the first and second errors are due to the uncertainty of the form factors and ξ, respectively.

IV. SUMMARY

We have studied the radiative and semileptonic B decays involving the tensor meson $K_2^*(1430)$ in the final states. Using the large energy effective theory techniques, $B \to K_2^*(1430)$ transition form factors have been formulated in the large recoil region. There are only two independent functions $\zeta_\perp(q^2)$ and $\zeta_\parallel(q^2)$ that describe all relevant form factors. We have determined the value of $\zeta_\perp(0)$ from the measurement of $\mathcal{B}(B^0 \to K_2^{*0}(1430)\gamma)$. Adopting a dipole q^2 dependency for the LEET functions and $\zeta_\parallel(q^2)/\zeta_\perp(q^2) = 1.0 \pm 0.2$, for which the former consists with the QCD counting rules, and the latter is favored by the $B \to \phi K_2^*$ data, we have investigated the decays $B \to K_2^*\ell^+\ell^-$ and $B \to K_2^*\nu\bar{\nu}$. Note that $\zeta_\parallel$ only gives corrections of order $m_{K_2^*}/m_B$. We have discussed two dedicated observables, the longitudinal distribution dF_L/ds and forward-backward asymmetry, in the $B \to K_2^*\ell^+\ell^-$ decay. Recent forward-backward asymmetry measurements for $B \to K^*\ell^+\ell^-$ decays [9,13,15] seem to (i) allow the possibility of flipping the sign of $c_7^{\rm eff}$, or (ii) have both c_9 and c_{10} flipped in sign, as compared with the SM. Meanwhile, in the large recoil region, *BABAR* has recently reported the large isospin asymmetry for the $B \to K^*\ell^+\ell^-$ decays, which qualitatively favors the flipped-sign $c_7^{\rm eff}$ model over the SM [22]. Therefore, in the present study, in addition to the SM, we focus the new physics effects on $c_7^{\rm eff}$ with the sign flipped. It should be note that the magnitude of $c_7^{\rm eff}$ is stringently constrained by the $B \to X_s\gamma$ data, which is consistent with the SM prediction.

For the $B \to K_2^*\ell^+\ell^-$ decay, of particular interest is the large recoil region, where the uncertainties of form factors are considerably reduced not only by taking the ratios of the form factors but also by computing in the large $E_{K_2^*}$ limit. In this region, where the invariant mass of the lepton pair $s \simeq 2-4$ GeV^2, due to the flipped sign of $c_7^{\rm eff}$ compared with the SM result, dF_L/ds is reduced by 20–30%, and its value can be ~0.8. One the other hand, in the SM the asymmetry zero is about 3.4 GeV^2, but changing the sign of $c_7^{\rm eff}$ yields a positive forward-backward asymmetry for all values of the invariant mass of the lepton pair.

We have obtained the branching fraction for $B \to K_2^*\nu\bar{\nu}$ in the SM. This mode enhanced by the summation over three light neutrinos is theoretically cleaner due to the absence of long-distance corrections related to the relevant four-fermion operators. This decay is relevant for the search for the nonstandard Z^0 coupling, light dark matter, and unparticles.

In summary, the investigation of the semileptonic B decays involving $K_2^*(1430)$ will further provide complementary information on physics beyond the standard model. Our results also exhibit the impressed resemblance of the physical properties between $B \to K_2^*(1430)\ell^+\ell^-$, $\nu\bar{\nu}$ and $B \to K^*(892)\ell^+\ell^-$, $\nu\bar{\nu}$.

ACKNOWLEDGMENTS

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[1] I. Adachi *et al.* (Belle Collaboration), arXiv:0810.0335.
[2] B. Aubert *et al.* (*BABAR* Collaboration), Phys. Rev. Lett. **102**, 091803 (2009).
[3] A. Yarritu *et al.* (*BABAR* Collaboration), arXiv:0808.1915.
[4] M. Nakao *et al.* (BELLE Collaboration), Phys. Rev. D **69**, 112001 (2004).
[5] T. E. Coan *et al.* (CLEO Collaboration), Phys. Rev. Lett. **84**, 5283 (2000).
[6] B. Aubert *et al.* (*BABAR* Collaboration), Phys. Rev. D **70**, 091105 (2004).
[7] S. Nishida *et al.* (Belle Collaboration), Phys. Rev. Lett. **89**, 231801 (2002).
[8] B. Aubert *et al.* (*BABAR* Collaboration), Phys. Rev. D **73**, 092001 (2006).
[9] I. Adachi (Belle Collaboration), arXiv:0810.0335.
[10] S. Anderson *et al.* (CLEO Collaboration), Phys. Rev. Lett. **87**, 181803 (2001).
[11] K. F. Chen *et al.* (BELLE Collaboration), Phys. Rev. Lett. **99**, 221802 (2007).
[12] B. Aubert *et al.* (*BABAR* Collaboration), Phys. Rev. D **78**, 072007 (2008).
[13] A. Ishikawa *et al.*, Phys. Rev. Lett. **96**, 251801 (2006).
[14] E. Barberio *et al.* (Heavy Flavor Averaging Group), arXiv:0808.1297.
[15] B. Aubert *et al.* (*BABAR* Collaboration), Phys. Rev. D **79**, 031102 (2009).
[16] W. J. Li, Y. B. Dai, and C. S. Huang, Eur. Phys. J. C **40**, 565 (2005).
[17] G. Burdman, Phys. Rev. D **52**, 6400 (1995).
[18] J. L. Hewett and J. D. Wells, Phys. Rev. D **55**, 5549 (1997).
[19] Y. G. Xu, R. M. Wang, and Y. D. Yang, Phys. Rev. D **74**, 114019 (2006).
[20] P. Colangelo, F. De Fazio, R. Ferrandes, and T. N. Pham, Phys. Rev. D **73**, 115006 (2006).
[21] C. H. Chen and C. Q. Geng, Phys. Rev. D **66**, 094018 (2002).
[22] G. Eigen, arXiv:0807.4076.
[23] T. Feldmann and J. Matias, J. High Energy Phys. 01 (2003) 074.
[24] A. Ali, P. Ball, L. T. Handoko, and G. Hiller, Phys. Rev. D **61**, 074024 (2000).
[25] H. Hatanaka and K. C. Yang, Phys. Rev. D **77**, 094023 (2008); **78**, 059902(E) (2008).
[26] H. Hatanaka and K. C. Yang, Phys. Rev. D **78**, 074007 (2008).
[27] H. Yang *et al.*, Phys. Rev. Lett. **94**, 111802 (2005).
[28] H. Y. Cheng and C. K. Chua, Phys. Rev. D **69**, 094007 (2004).
[29] D. Ebert, R. N. Faustov, and V. O. Galkin, Phys. Rev. D **64**, 094022 (2001); D. Ebert, R. N. Faustov, V. O. Galkin, and H. Toki, Phys. Rev. D **64**, 054001 (2001).
[30] A. Datta, Y. Gao, A. V. Gritsan, D. London, M. Nagashima, and A. Szynkman, Phys. Rev. D **77**, 114025 (2008).
[31] J. Charles, A. Le Yaouanc, L. Oliver, O. Pene, and J. C. Raynal, Phys. Rev. D **60**, 014001 (1999).
[32] G. Buchalla, G. Hiller, and G. Isidori, Phys. Rev. D **63**, 014015 (2000).
[33] C. Bird, P. Jackson, R. V. Kowalewski, and M. Pospelov, Phys. Rev. Lett. **93**, 201803 (2004).
[34] H. Georgi, Phys. Rev. Lett. **98**, 221601 (2007).
[35] T. M. Aliev, A. S. Cornell, and N. Gaur, J. High Energy Phys. 07 (2007) 072.
[36] E. R. Berger, A. Donnachie, H. G. Dosch, and O. Nachtmann, Eur. Phys. J. C **14**, 673 (2000).
[37] C. Amsler *et al.* (Particle Data Group), Phys. Lett. B **667**, 1 (2008).
[38] CKMfitter Group (http://ckmfitter.in2p3.fr), results as of 2008 summer.
[39] M. Beneke, T. Feldmann, and D. Seidel, Nucl. Phys. **B612**, 25 (2001).
[40] A. Ali and A. Y. Parkhomenko, Eur. Phys. J. C **23**, 89 (2002).
[41] S. W. Bosch and G. Buchalla, Nucl. Phys. **B621**, 459 (2002).
[42] V. L. Chernyak and A. R. Zhitnitsky, Phys. Rep. **112**, 173 (1984).
[43] E. Barberio *et al.* (Heavy Flavor Averaging Group), arXiv:0704.3575 and online update at http://www.slac.-stanford.edu/xorg/hfag.
[44] C. S. Lim, T. Morozumi, and A. I. Sanda, Phys. Lett. B **218**, 343 (1989).
[45] A. Ali, T. Mannel, and T. Morozumi, Phys. Lett. B **273**, 505 (1991).
[46] J. Matias, Nucl. Phys. B, Proc. Suppl. **185**, 68 (2008).
[47] C. Bobeth, G. Hiller, and G. Piranishvili, J. High Energy Phys. 07 (2008) 106.
[48] A. J. Buras and R. Fleischer, Adv. Ser. Dir. High Energy Phys. **15**, 65 (1998).
[49] A. Ali and T. Mannel, Phys. Lett. B **264**, 447 (1991); **274**, 526(E) (1992).
[50] T. Inami and C. S. Lim, Prog. Theor. Phys. **65**, 297 (1981); **65**, 1772(E) (1981).
[51] G. Buchalla and A. J. Buras, Nucl. Phys. **B400**, 225 (1993).
[52] G. Buchalla, A. J. Buras, and M. E. Lautenbacher, Rev. Mod. Phys. **68**, 1125 (1996).
[53] G. Buchalla and A. J. Buras, Nucl. Phys. **B548**, 309 (1999).
[54] P. Colangelo, F. De Fazio, P. Santorelli, and E. Scrimieri, Phys. Lett. B **395**, 339 (1997).
[55] D. Melikhov, N. Nikitin, and S. Simula, Phys. Lett. B **428**, 171 (1998).
[56] C. S. Kim, Y. G. Kim, and T. Morozumi, Phys. Rev. D **60**, 094007 (1999).