Study of $B \to K_0^*(1430) l\bar{l}$ decays

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We examine the exclusive rare decays of $B \to K_0^*(1430)\ell\bar{\ell}$ with $K_0^*(1430)$ being the p-wave scalar meson and $\ell = \nu, e, \mu, \tau$ in the standard model. The form factors for the $B \to K_0^*$ transition matrix elements are evaluated in the light-front quark model. For the decays of $B \to K_0^*\ell\bar{\ell}$, the branching ratios are found to be $(11.6, 1.63, 1.62, 0.029) \times 10^{-7}$ with $\ell = (\nu, e, \mu, \tau)$ and the integrated longitudinal lepton polarization asymmetries $(-0.97, -0.95, -0.03)$ with $\ell = (e, \mu, \tau)$, respectively.

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I. INTRODUCTION

The suppressed inclusive flavor-changing neutral current (FCNC) process of $B \to X_s\ell^+\ell^-$, induced by electroweak penguin and box diagrams in the standard model (SM), has been observed by both *BABAR* [1] and Belle [2] with the branching ratio (BR) of $(4.5 \pm 1.0) \times 10^{-6}$ [3] for dilepton masses greater than 0.2 GeV, where ℓ is either an electron or a muon and X_s is a hadronic recoil system that contains a kaon. The exclusive decays of $B \to K\ell^+\ell^-$ and $B \to K^*(892)\ell^+\ell^-$ have also been measured with the BRs [3] of $(0.54 \pm 0.08) \times 10^{-6}$ and $(1.05 \pm 0.20) \times 10^{-6}$ [4,5], which agree with the theoretically estimated values [6–9], respectively.

There have been many investigations of rare B semileptonic decays of induced by the FCNC transition of $b \to s$ [10] since the CLEO observation [11] of $b \to s\gamma$. The studies are even more complete if similar studies for the p-wave mesons of B decays such as $B \to K_{0,2}^*(1430)\ell\bar{\ell}$ and $B \to K_{1A,1B}\ell\bar{\ell}$ are also included. In fact, the study of $B \to K_2^*(1430)\ell^+\bar{\ell}^-$ has been done in Ref. [12]. It is clear that these FCNC rare decays are important for not only testing the SM but probing new physics. In this report, we concentrate on the exclusive rare decays of $B \to K_0^*\ell\bar{\ell}$, where K_0^* represents the p-wave scalar meson of $K_0^*(1430)$ and ℓ stands for a charged lepton or neutrino. To obtain the decay rates and branching ratios, we need to calculate the transition form factors of $B \to K_0^*$ due to the axial-vector and axial-tensor currents, respectively, in the standard model. We will use the framework of the light-front quark model (LFQM) [13–15] to evaluate these form factors.

This report is organized as follows: We present the relevant formulas in Sec. II. First, we give the effective Hamiltonians for $B \to K_0^*\ell\bar{\ell}$ induced by $b \to s\ell\bar{\ell}$. Then, we calculate the hadronic form factors for the $B \to K_0^*$ transition in the LFQM. Finally, we study the branching ratios and polarization asymmetries of the decays. In Sec. III, we show our numerical results on form factors and the physical quantities of the decays. We give our conclusions in Sec. IV.

II. THE FORMULAS

To study the exclusive decays of $B \to K_0^*\ell\bar{\ell}$, we start with the effective Hamiltonians at the quark level, given by

$$
\begin{aligned}
\mathcal{H}(b \to s\nu\bar{\nu}) &= \frac{G_F}{\sqrt{2}}\frac{\alpha}{2\pi \sin^2\theta_W}\lambda_t D(x_t)\bar{b}\gamma_\mu(1-\gamma_5) \\
&\quad \times s\bar{\nu}_\ell\gamma_\mu(1-\gamma_5)\nu_\ell, \\
\mathcal{H}(b \to s\ell^+\ell^-) &= \frac{G_F\alpha\lambda_t}{\sqrt{2}\pi}\bigg[C_9^{\mathrm{eff}}(m_b)\bar{s}_L\gamma_\mu b_L\bar{\ell}\gamma^\mu\ell \\
&\quad + C_{10}\bar{s}_L\gamma_\mu b_L\bar{\ell}\gamma^\mu\gamma_5\ell \\
&\quad - \frac{2m_bC_7(m_b)}{q^2}\bar{s}_L i\sigma_{\mu\nu}q^\nu b_R\bar{\ell}\gamma^\mu\ell\bigg],
\end{aligned} \tag{1}
$$

where $x_t \equiv m_t^2/m_W^2$, $\lambda_t = V_{ts}^*V_{tb}$, $D(x_t)$ is the top-quark loop function [16,17] and C_i are the Wilson coefficients (WCs) with their explicit expressions given in Ref. [18]. In particular, C_9^{eff}, which contains the contribution from the on-shell charm-loop, is given by [18]

$$
\begin{aligned}
C_9^{\mathrm{eff}}(\mu) &= C_9(\mu) + (3C_1(\mu) + C_2(\mu))h(z,\hat{s}), \\
h(z,\hat{s}) &= -\frac{8}{9}\ln\frac{m_b}{\mu} - \frac{8}{9}\ln z + \frac{8}{27} + \frac{4}{9}x - \frac{2}{9}(2+x)|1-x|^{1/2} \\
&\quad \times \begin{cases} \ln\left|\frac{\sqrt{1-x}+1}{\sqrt{1-x}-1}\right| - i\pi, & \text{for} x \equiv 4z^2/\hat{s} < 1, \\ 2\arctan\frac{1}{\sqrt{x-1}}, & \text{for} x \equiv 4z^2/\hat{s} > 1, \end{cases}
\end{aligned} \tag{2}
$$

where $z = m_c/m_b$ and $\hat{s} = q^2/m_b^2$ with q^2 being the invariant mass of the dilepton. Here, we have ignored the resonant contributions [19,20] as the modes such as $B \to J/\Psi K_0^*$ have not be seen yet. To calculate the decay rates, we need to evaluate the hadronic matrix elements for the $B \to K_0^*$ transition involving the axial-vector and axial-tensor currents, whereas those from the vector and tensor ones are zero due to the parity conservation in strong interactions. In the following, we will study the matrix elements in the LFQM [21,22], which has been successfully applied to many weak processes with the heavy-to-heavy and heavy-to-light transitions in the timelike regions [13–15,23].

In the LFQM, a meson bound state consisting of a heavy quark q_1 and an antiquark q_2 with the total momentum P and spin S can be written as

$$|B(P, S, S_z)\rangle = \int \frac{dp_1^+ d^2 p_{1\perp}}{2(2\pi)^3} \frac{dp_2^+ d^2 p_{2\perp}}{2(2\pi)^3} \times 2(2\pi)^3 \delta^3(\tilde{P} - \tilde{p}_1 - \tilde{p}_2) \times \sum_{\lambda_1, \lambda_2} \Psi^{SS_z}(\tilde{p}_1, \tilde{p}_2, \lambda_1, \lambda_2)|q_1(p_1, \lambda_1) \times \bar{q}_2(p_2, \lambda_2)\rangle, \tag{3}$$

where p_1 and p_2 are the on-mass-shell light-front momenta

$$\tilde{p} = (p^+, p_\perp), \qquad p_\perp = (p^1, p^2), \qquad p^- = \frac{m^2 + p_\perp^2}{p^+}, \tag{4}$$

with

$$p_1^+ = (1 - x)P^+, \qquad p_2^+ = xP^+, \qquad p_{1\perp} = (1 - x)P_\perp + k_\perp, \qquad p_{2\perp} = xP_\perp - k_\perp. \tag{5}$$

Here $(x, k_\perp)$ are the light-front relative momentum variables and $\vec{k}_\perp$ is the component of the internal momentum $\vec{k} = (\vec{k}_\perp, k_z)$. The momentum-space wave-function Ψ^{SS_z} in Eq. (3) can be expressed as

$$\Psi^{SS_z}(\tilde{p}_1, \tilde{p}_2, \lambda_1, \lambda_2) = R^{SS_z}_{\lambda_1 \lambda_2}(x, k_\perp)\phi^{(p)}(x, k_\perp), \tag{6}$$

where $R^{SS_z}_{\lambda_1 \lambda_2}$ constructs a state of a definite spin (S, S_z) out of the light-front helicity (λ_1, λ_2) eigenstates and $\phi^{(p)}(x, k_\perp)$ describes the momentum distribution of the constituents in the bound state for the s-wave (p-wave) meson. Explicitly, it is more convenient to use the covariant form for $R^{SS_z}_{\lambda_1 \lambda_2}$ [15], given by

$$R^{SS_z}_{\lambda_1 \lambda_2}(x, k_\perp) = h_M \bar{u}(p_1, \lambda_1)\Gamma_M v(p_2, \lambda_2), \qquad (M = P_i, S_f), \tag{7}$$

where

$$\Gamma_{P_i} = \gamma_5, h_{P_i} = (m_{P_i}^2 - M_0^2)\sqrt{\frac{x(1 - x)}{N_c}} \frac{1}{\sqrt{2}5_0; M}, \qquad \Gamma_{S_f} = -i, \qquad h_{S_f} = (m_{S_f}^2 - M_0^2)\sqrt{\frac{x(1 - x)}{N_c}} \frac{1}{\sqrt{2}5_0; M} \frac{5_0^2; M}{\sqrt{3}M_0}, \tag{8}$$

for the initial s-wave pseudoscalar (P_i) and final p-wave scalar (S_f) mesons of B and K_0^*, respectively, with

$$\tilde{M}_0 \equiv \sqrt{M_0^2 - (m_1 - m_2)^2}, \qquad M_0^2 = \frac{m_1^2 + k_\perp^2}{(1 - x)} + \frac{m_2^2 + k_\perp^2}{x}. \tag{9}$$

In our calculations, we use the Gaussian type wave functions

$$\phi(x, k_\perp) = 4\left(\frac{\pi}{\omega_M^2}\right)^{3/4} \sqrt{\frac{dk_z}{dx}} \exp\left(-\frac{\vec{k}^2}{2\omega_M^2}\right), \qquad \phi^p(x, k_\perp) = \sqrt{\frac{2}{\omega_M^2}}\phi(x, k_\perp), \tag{10}$$

for the s-wave and p-wave mesons, respectively, where ω_M is the meson scale parameter and k_z is defined through

$$1 - x = \frac{e_1 - k_z}{e_1 + e_2}, \qquad x = \frac{e_2 + k_z}{e_1 + e_2}, \tag{11}$$

with $e_i = \sqrt{m_i^2 + \vec{k}^2}$. Then, we have

$$M_0 = e_1 + e_2, \qquad k_z = \frac{xM_0}{2} - \frac{m_2^2 + k_\perp^2}{2xM_0}, \tag{12}$$

and

$$\frac{dk_z}{dx} = \frac{e_1 e_2}{x(1 - x)M_0}. \tag{13}$$

We normalize the meson state as

$$\langle B(P', S', S_z')|B(P, S, S_z)\rangle = 2(2\pi)^3 P^+ \delta^3(\tilde{P}' - \tilde{P})\delta_{S'S}\delta_{S_z'S_z}, \tag{14}$$

so that the normalization condition of the momentum distribution function can be obtained by

$$\int \frac{dx d^2 k_\perp}{2(2\pi)^3} |\phi(x, k_\perp)|^2 = 1. \tag{15}$$

We are now ready to calculate the matrix elements of the $P_i \to S_f$ transition, which can be defined by

$$\langle S_f(p_f)|A_\mu|P_i(p_i)\rangle = -i[u_+(q^2)P_\mu + u_-(q^2)q_\mu], \qquad \langle S_f(p_f)|T^5_{\mu\nu}q^\nu|P_i(p_i)\rangle = \frac{-i}{m_{P_i} + m_{S_f}} \times [q^2 P_\mu - (P \cdot q)q_\mu]F_T(q^2), \tag{16}$$

where $A_\mu = \bar{q}_f \gamma_\mu \gamma_5 q_i$, $T^5_\mu = \bar{q}_f i\sigma_{\mu\nu}\gamma_5 q_i$, $P = p_i + p_f$ and $q = p_i - p_f$ with the initial (final) meson bound state $q_i\bar{q}_3$ $(q_f\bar{q}_3)$. We note that all form factors will be studied in the timelike physical meson decay region of $0 \le q^2 \le (m_{P_i} - m_{S_f})^2$. The form factors in Eq. (16) are found to be

$$u_+(q^2) = \frac{(1-r_-)H(r_+) - (1-r_+)H(r_-)}{r_+ - r_-},$$
$$u_-(q^2) = \frac{(1+r_-)H(r_+) - (1+r_+)H(r_-)}{r_+ - r_-}, \tag{17}$$
$$F_T(q^2) = -\int_0^r dx \int \frac{d^2k_\perp}{2(2\pi)^3} \frac{\widetilde{M}_0^2}{2\sqrt{3}M_0} \phi_{S_f}^{p*}(x', k_\perp)\phi_{P_i}(x, k_\perp) \frac{m_{P_i} + m_{S_f}}{(1+2r)q^2 - (m_{P_i}^2 - m_{S_f}^2)} \frac{A}{\sqrt{\mathcal{A}_{P_i}^2 + k_\perp^2}\sqrt{\mathcal{A}_{S_f}^2 + k_\perp^2}},$$

where $r \equiv p_f^+/p_i^+$, $x' = x/r$,

$$r_\pm = \frac{m_{S_f}}{m_{P_i}}\left[v_i \cdot v_f \pm \sqrt{(v_i \cdot v_f)^2 - 1}\right], \qquad \left(v_i \cdot v_f = \frac{m_{P_i}^2 + m_{S_f}^2 - q^2}{2m_{P_i}m_{S_f}}\right),$$
$$H(r) = -\int_0^r dx \int \frac{d^2k_\perp}{2(2\pi)^3} \frac{\widetilde{M}_0^2}{2\sqrt{3}M_0} \phi_{S_f}^{p*}(x', k_\perp)\phi_{P_i}(x, k_\perp) \times \frac{[m_{q_i}x + m_{q_3}(1-x)][-m_{q_f}x' + m_{q_3}(1-x')] + k_\perp^2}{\sqrt{\mathcal{A}_{P_i}^2 + k_\perp^2}\sqrt{\mathcal{A}_{S_f}^2 + k_\perp^2}},$$
$$A = \frac{1}{\sqrt{xx'}(1-x)(1-x')}\{[xm_{q_i} + (1-x)m_{q_3}][-x'm_{q_f} + (1-x')m_{q_3}][x(1-x')m_{q_i} - x'(1-x)m_{q_3}]$$
$$+ k_\perp^2[-x'(1-x')(2x-1)m_{q_f} + (x-x')(x+x'-2xx')m_{q_3} + x(1-x)(1-2x')m_{q_i}]\}, \tag{18}$$
$$\mathcal{A}_{P_i(S_f)} = m_{q_i(q_f)}x^{(\prime)} + m_{q_3}(1 - x^{(\prime)}).$$

The sign $+(-)$ of $r_{+(-)}$ represents the final meson recoiling in the positive (negative) z direction relative to the initial meson.

We note that to evaluate the form factors, we have to fix the meson scale parameters ω_B and $\omega_{K_0^*}$ in the meson wave functions in Eq. (10) by some known parameters such as the meson decay constants, defined by

$$f_{P_i} = \sqrt{24}\int \frac{dxd^2k_\perp}{2(2\pi)^3} \phi(x, k_\perp) \frac{m_{q_i}x + m_{q_3}(1-x)}{\sqrt{\mathcal{A}_{P_i}^2 + k_\perp^2}}, \qquad f_{S_f} = \sqrt{24}\int \frac{dx'd^2k_\perp}{2(2\pi)^3} \frac{\widetilde{M}_0^2}{2\sqrt{3}M_0} \phi^p(x', k_\perp) \frac{m_{q_f}x' - m_{q_3}(1-x')}{\sqrt{\mathcal{A}_{S_f}^2 + k_\perp^2}}. \tag{19}$$

By using Eqs. (1) and (16), we derive the differential decay rates of $B \to K_0^* \ell\bar{\ell}$ as

$$\frac{d\Gamma(B \to K_0^*\nu\bar{\nu})}{ds} = \frac{G_F^2|\lambda_t|^2\alpha_{em}^2|D(x_t)|^2 m_B^5}{2^8\pi^5\sin^4\theta_W}\varphi_{K_0^*}^{1/2}|u_+|^2,$$
$$\frac{d\Gamma(B \to K_0^*\ell^+\ell^-)}{ds} = \frac{G_F^2|\lambda_t|^2 m_B^5\alpha_{em}^2}{3\cdot 2^9\pi^5}\left(1 - \frac{4t}{s}\right)^{1/2}\varphi_{K_0^*}^{1/2}\left[\left(1 + \frac{2t}{s}\right)\alpha_{K_0^*} + t\delta_{K_0^*}\right], \tag{20}$$

where

$$s = q^2/m_B^2, \qquad t = m_l^2/m_B^2, \qquad r_{K_0^*} = m_{K_0^*}^2/m_B^2, \qquad \varphi_{K_0^*} = (1 - r_{K_0^*})^2 - 2s(1 + r_{K_0^*}) + s^2,$$
$$\alpha_{K_0^*} = \varphi_{K_0^*}\left(|C_9^{\rm eff}u_+ - \frac{2C_7F_T}{1 + \sqrt{r_{K_0^*}}}|^2 + |C_{10}u_+|^2\right), \tag{21}$$
$$\delta_{K_0^*} = 6|C_{10}|^2\{[2(1 + r_{K_0^*}) - s]|u_+|^2 + 2(1 - r_{K_0^*})\,{\rm Re}(u_+u_-^*) + s|u_-|^2\}.$$

When the polarization of the charged lepton is in the longitudinal direction, i.e. $\hat{n} = \mathbf{e}_L = \vec{p}_\ell/|\vec{p}_\ell| = \pm 1$, we can also define the longitudinal lepton polarization asymmetry in $B \to K_0^*\ell^+\ell^-$ as follows [24,25]:

$$P_L(s) = \frac{\frac{d\Gamma(\hat{n}=-1)}{ds} - \frac{d\Gamma(\hat{n}=1)}{ds}}{\frac{d\Gamma(\hat{n}=-1)}{ds} + \frac{d\Gamma(\hat{n}=1)}{ds}}. \tag{22}$$

From Eq. (22), we find that

$$P_L = \frac{2(1 - \frac{4t}{s})^{1/2}}{(1 + \frac{2t}{s})\alpha_{K_0^*} + t\delta_{K_0^*}} \times {\rm Re}\left[\varphi_{K_0^*}\left(C_9^{\rm eff}u_+ - 2\frac{C_7F_T}{1 + \sqrt{r_{K_0^*}}}\right)(C_{10}u_+)^*\right] \tag{23}$$

in $B \to K_0^*\ell^+\ell^-$. We remark that there is no forwardbackward asymmetry for $B \to K_0^* \ell^+ \ell^-$ in the SM similar to other pseudoscalar to pseudoscalar dilepton decays such as $P_i \to P_f \ell^+ \ell^-$ with $P_i = K(B)$ and $P_f = \pi(K)$ [26].

III. NUMERICAL RESULTS

In our numerical study of the hadronic matrix elements for the $B \to K_0^*$ transition, we fix the quark masses to be $m_b = 4.64$, $m_s = 0.37$, and $m_{u,d} = 0.26$ GeV and use the meson decay constants to determine the meson scale parameters as shown in Table I. We note that the current direct measurement of f_B is $0.229^{+0.036+0.034}_{-0.031-0.037}$ GeV [27], whereas there is no experimental information on $f_{K_0^*}$. Since f_B and $f_{K_0^*}$ are fixed to be 0.18 and 0.021 GeV in Ref. [15], respectively, we will also use these values in our numerical analysis and briefly discuss other values at the end. Our results for the form factors at $q^2 = 0$ are given in Table II. In the table, as a comparison, we have also shown the value of $u_+(0)$ used in Ref. [15]. To give explicit q^2 dependent form factors, we fit our results to the form

$$F(q^2) = \frac{F(0)}{1 - a(q^2/m_B^2) + b(q^2/m_B^2)^2}, \qquad (24)$$

with the fitted ranges of $F_T(q^2)$ and $u_\pm(q^2)$ being $0 \le q^2 \le 12$ and $0 \le q^2 \le (m_B - m_{K_0^*})^2$ GeV^2, respectively. In Fig. 1, the form factors as functions of q^2 are presented, where (a) $u_\pm(q^2)$ and (b) $F_T(q^2)$. In Table III, we give the values of the relevant WCs at the scale of $\mu \sim 4.8$ GeV [9]. With $|\lambda_t| \simeq 0.041$, we illustrate the differential decay branching ratios for $B \to K_0^* \nu \bar{\nu}$ and $B \to K_0^* \ell^+ \ell^- (\ell = \mu, \tau)$ as functions of s in Figs. 2 and 3, respectively. By integrating the differential ratios over $s = q^2/m_B^2$ for $B \to K_0^* \nu \bar{\nu}$ and $B \to K_0^* \ell^+ \ell^-$, we obtain

$$Br(B \to K_0^* \nu \bar{\nu}) = 1.16 \times 10^{-6} \qquad (25)$$

and

$$Br(B \to K_0^* e^+ e^-, K_0^* \mu^+ \mu^-, K_0^* \tau^+ \tau^-) = 1.63 \times 10^{-7}, \quad 1.62 \times 10^{-7}, \quad 2.86 \times 10^{-9}, \qquad (26)$$

respectively. Note that the small branching ratio of the tau mode is due to the highly suppressed phase space as shown in Fig. 3. We also present the longitudinal lepton polarization asymmetries of $B \to K_0^* \ell^+ \ell^-$ as functions of s in Fig. 4. We note that our results for the electron mode are

TABLE I. Meson decay constants and scale parameters (in units of GeV).

f_B	ω_B	$f_{K_0^*}$	$\omega_{K_0^*}$
0.16	0.4763	0.015	0.2106
0.18	0.5239	0.021	0.3001
0.20	0.5713	0.025	0.3837

TABLE II. Form factors at $q^2 = 0$ for the $B \to K_0^*$ transition.

	This work			[15]		
	$F(0)$	a	b	$F(0)$	a	b
u_+	−0.26	1.36	0.86	−0.26	1.52	0.64
u_-	0.21	1.26	0.93	...	...	...
F_T	0.34	1.64	1.72	...	...	...

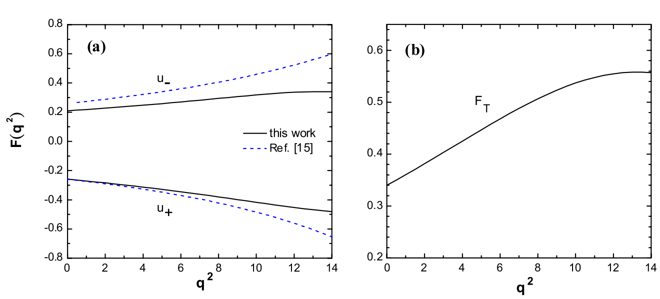

FIG. 1 (color online). Form Factors as functions of q^2 for (a) $u_+(q^2)$ and $u_-(q^2)$ and (b) $F_T(q^2)$.

TABLE III. Wilson coefficients for $m_t = 170$ GeV and $\mu = 4.8$ GeV.

WC	C_1	C_2	C_7	C_9	C_{10}
	−0.226	1.096	−0.305	4.186	−4.559

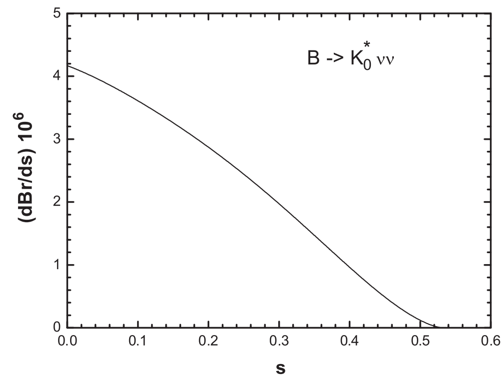

FIG. 2. Differential decay branching ratio for $B \to K_0^* \nu \bar{\nu}$ as a function of $s = q^2/m_B^2$.

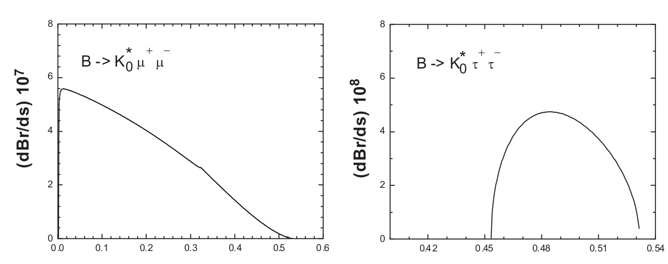

FIG. 3. Differential decay branching ratios for $B \to K_0^* \mu^+ \mu^-$ and $B \to K_0^* \tau^+ \tau^-$ as functions of $s = q^2/m_B^2$.

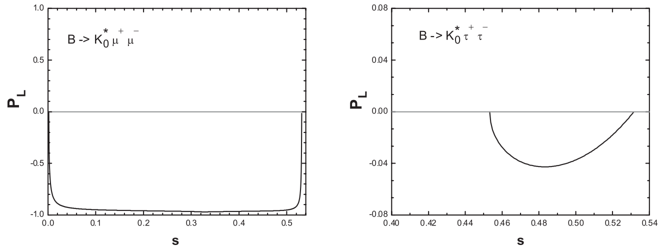

FIG. 4. Longitudinal lepton polarization asymmetries for $B \to K_0^* \ell^+ \ell^-$.

similar to those for the muon one. As shown in Fig. 4, $P_L(s)[B \to K_0^* \mu^+ \mu^-]$ is close to -1 except those close to the end points of $q^2_{\min} = 4m^2_\mu$ and $q^2_{\max} = (m_B - m_{K_0^*})^2$ at which they are zero and $P_L(s)[B \to K_0^* \tau^+ \tau^-]$ ranges from -0.5 to 0, while the integrated values of P_L are -0.97, -0.95, and -0.03 for electron, muon and tau modes, respectively. It is clear that due to the efficiency for the detectability of the tau lepton, it is impossible to measure the tau lepton polarization in the near future.

Finally, we remark that our results are insensitive (sensitive) to the value of f_B ($f_{K_0^*}$). For examples, $Br(B \to K_0^* \ell \bar{\ell})$ ($\ell = \nu, e, \mu$) decrease about 6% by increasing $f_B =$ 0.16 to 0.20 GeV, while they increase about 50% by increasing $f_{K_0^*} = 0.021$ to 0.025 GeV.

IV. CONCLUSIONS

We have studied the exclusive rare decays of $B \to K_0^* \ell \bar{\ell}$. We have calculated the form factors for the $B \to K_0^*$ transition matrix elements in the LFQM. We have evaluated the decay branching ratios and the longitudinal charged-lepton polarization asymmetries in the SM. Explicitly, we have found that $Br(B \to K_0^* \ell \bar{\ell})(\ell = \nu, e, \mu, \tau) = (11.6, 1.63, 1.62, 0.029) \times 10^{-7}$ and the integrated longitudinal lepton polarization asymmetries of $B \to K_0^* \ell^+ \ell^- (\ell = e, \mu, \tau)$ are -0.97, -0.95 and -0.03, respectively. It is clear that some of the above p-wave B decays and asymmetries can be measured at the ongoing as well as future B factories.

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[1] B. Aubert *et al.* (*BABAR* Collaboration), Phys. Rev. Lett. **93**, 081802 (2004).
[2] M. Iwasaki *et al.* (Belle Collaboration), Phys. Rev. D **72**, 092005 (2005).
[3] W. M. Yao *et al.* (Particle Data Group), J. Phys. G **33**, 1 (2006).
[4] B. Aubert *et al.* (*BABAR* Collaboration), Phys. Rev. Lett. **91**, 221802 (2003).
[5] A. Ishikawa *et al.* (Belle Collaboration), Phys. Rev. Lett. **91**, 261601 (2003).
[6] A. Ali, P. Ball, L. T. Handoko, and G. Hiller, Phys. Rev. D **61**, 074024 (2000); P. Colangelo, F. De Fazio, P. Santorelli, and E. Scrimieri, Phys. Rev. D **53**, 3672 (1996); D. Melikhov, N. Nikitin, and S. Simula, Phys. Rev. D **57**, 6814 (1998).
[7] C. Q. Geng and C. P. Kao, Phys. Rev. D **57**, 4479 (1998).
[8] A. Ali, E. Lunghi, C. Greub, and G. Hiller, Phys. Rev. D **66**, 034002 (2002); T. M. Aliev, H. Koru, A. Özpineci, and M. Savci, Phys. Lett. B **400**, 194 (1997); T. M. Aliev, M. Savci, and A. Özpineci, Phys. Rev. D **56**, 4260 (1997); G. Burdman, Phys. Rev. D **52**, 6400 (1995); N. G. Deshpande and J. Trampetic, Phys. Rev. Lett. **60**, 2583 (1988); C. Greub, A. Ioannissian, and D. Wyler, Phys. Lett. B **346**, 149 (1995); J. L. Hewett and J. D. Wells, Phys. Rev. D **55**, 5549 (1997); and references therein.
[9] C. H. Chen and C. Q. Geng, Phys. Rev. D **63**, 114025 (2001); **64**, 074001 (2001).
[10] For a recent review, see A. Ali *et al.*, Phys. Rev. D **61**, 074024 (2000).
[11] M. S. Alam *et al.* (CLEO Collaboration), Phys. Rev. Lett. **74**, 2885 (1995).
[12] S. Rai Choudhury, A. S. Cornell, G. C. Joshi, and B. H. J. McKellar, Phys. Rev. D **74**, 054031 (2006).
[13] H. Y. Cheng, C. Y. Cheung, and C. W. Hwang, Phys. Rev. D **55**, 1559 (1997).
[14] C. Q. Geng, C. W. Hwang, C. C. Lih, and W. M. Zhang, Phys. Rev. D **64**, 114024 (2001).
[15] H. Y. Cheng, C. K. Chua, and C. W. Hwang, Phys. Rev. D **69**, 074025 (2004).
[16] G. Belanger and C. Q. Geng, Phys. Rev. D **43**, 140 (1991).
[17] G. Guchalla and A. J. Buras, Nucl. Phys. **B400**, 225 (1993).
[18] G. Buchalla, A. J. Buras, and M. E. Lautenbacher, Rev. Mod. Phys. **68**, 1125 (1996).
[19] C. S. Lim, T. Morozumi, and A. I. Sanda, Phys. Lett. B **218**, 343 (1989); N. G. Deshpande, J. Trampetic, and K. Panose, Phys. Rev. D **39**, 1461 (1989); P. J. O'Donnell and H. K. K. Tung, Phys. Rev. D **43**, R2067 (1991).
[20] C. H. Chen and C. Q. Geng, Phys. Rev. D **66**, 034006 (2002).
[21] M. V. Terent'ev, Sov. J. Nucl. Phys. **24**, 106 (1976); V. B. Berestetsky and M. V. Terent'ev, *ibid.* **24**, 547 (1976); **25**, 347 (1977).
[22] P. L. Chung, F. Coester, and W. N. Polyzou, Phys. Lett. B **205**, 545 (1988).

[23] C. Y. Cheung, C. W. Hwang, and W. M. Zhang, Z. Phys. C **75**, 657 (1997); C. Q. Geng, C. C. Lih, and Wei-Min Zhang, Phys. Rev. D **57**, 5697 (1998); N. B. Demchuk, I. L. Grach, I. M. Narodetskii, and S. Simula, Phys. At. Nucl. **59**, 2152 (1996) [Yad. Fiz. **59**, 2235 (1996)]; I. L. Grach, I. M. Narodetsky, and S. Simula, Phys. Lett. B **385**, 317 (1996).
[24] G. Belanger, C. Q. Geng, and P. Turcotte, Nucl. Phys. **B390**, 253 (1993).
[25] C. Q. Geng and C. P. Kao, Phys. Rev. D **54**, 5636 (1996).
[26] C. H. Chen, C. Q. Geng, and I. L. Ho, Phys. Rev. D **67**, 074029 (2003); C. H. Chen, C. Q. Geng, and A. K. Giri, Phys. Lett. B **621**, 253 (2005).
[27] K. Ikado *et al.*, Phys. Rev. Lett. **97**, 251802 (2006).