Lepton polarization asymmetry in $B \to K^{(*)} l^+ l^-$

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(Received 8 July 1996)

We investigate the longitudinal lepton polarization asymmetry in the exclusive processes $B \to K l^+ l^-$ and $B \to K^* l^+ l^-$. We include both short- and long-distance contributions to the asymmetry in our discussions. We find that average values of the polarization asymmetries of the muon and τ for $B \to K^{(*)} \mu^+ \mu^-$ and $B \to K^{(*)} \tau^+ \tau^-$ are -0.8 (-0.7) and -0.2 (-0.5), respectively. [S0556-2821(96)05321-0]

PACS number(s): 13.20.He, 13.88.+e

I. INTRODUCTION

As is well known, the study of B-meson physics is important in the determination of the elements of the Cabibbo-Kobayashi-Maskawa (CKM) [1] matrix of charged-current weak couplings. In addition, it could shed light on physics beyond the standard model. Recently, much interest [3] has been centered on rare B-meson decays induced by the flavor-changing neutral current $b \to s$ transition due to the CLEO measurement of the radiative $b \to s \gamma$ decay [2]. In the standard model, these rare processes do not occur at tree level and appear only at the quantum (loop) level. The short-distance (SD) contributions to the decays involving the $b \to s$ transition are dominated by loops with the top quark and they are basically free of the uncertainties in the CKM parameters. The rare B-meson decays are, thus, a good probe of heavy top quark physics as well as physics beyond the standard model.

It has been pointed out by Ali *et al.* [4] that, in the standard model, the measurement of the forward-backward asymmetry of the dileptons in the inclusive decays $b \to s l^+ l^-$ provides information on the short-distance contributions dominated by the top quark loops. Recently, it has been emphasized by Hewett [5] that the longitudinal lepton polarizations, which are other parity-violating observables, are also important asymmetries. In particular, the τ polarization in $b \to s \tau^+ \tau^-$ mode could be accessible to the B factories currently under construction. It is interesting to note that the dilepton forward-backward asymmetry of the exclusive decays $B \to M l^+ l^-$ is identically zero when M are pseudoscalar mesons such as π and K, but nonzero when M are vector mesons such as ρ and K^*. However, the longitudinal lepton polarizations for the exclusive modes are nonzero for both cases of pseudoscalar and vector mesons. In this paper we will examine these lepton polarization asymmetries for the exclusive decays of $B \to K l^+ l^-$ and $B \to K^* l^+ l^-$, in which we only concentrate on the muon and τ dileptonic modes since the electron polarization is hard to be measured experimentally. In our discussions, we will use the results of the relativistic quark model by the light-front formalism [6,7] on the form factors in the hadronic matrix elements between the B meson and the kaons.

The paper is organized as follows. In Sec. II, we give the general effective Hamiltonian in the standard model for the rare dilepton decays of interest. We study the exclusive decays of $B \to K l^+ l^-$ and $B \to K^* l^+ l^-$ in Sec. III and Sec. IV, respectively. Our conclusions are summarized in Sec. V.

II. EFFECTIVE HAMILTONIAN

The effective Hamiltonian relevant to $b \to s l^+ l^-$ based on an operator product expansion is given by [8]

$$\mathcal{H}_{\rm eff} = \frac{4 G_F \lambda_t}{\sqrt{2}} \sum_{i=1}^{10} C_i(\mu) O_i(\mu), \qquad (1)$$

where G_F denotes the Fermi constant, $\lambda_t = V_{tb} V_{ts}^*$ are the products of the CKM matrix elements. The $O_i(\mu)$ are the operators

$$O_1 = (\overline{s_\alpha} \gamma^\mu L b_\alpha)(\overline{c_\beta} \gamma_\mu L c_\beta),$$

$$O_2 = (\overline{s_\alpha} \gamma^\mu L b_\beta)(\overline{c_\beta} \gamma_\mu L c_\alpha),$$

$$O_3 = (\overline{s_\alpha} \gamma^\mu L b_\alpha)[(\overline{u_\beta} \gamma_\mu L u_\beta) + \cdots + (\overline{b_\beta} \gamma_\mu L b_\beta)],$$

$$O_4 = (\overline{s_\alpha} \gamma^\mu L b_\beta)[(\overline{u_\beta} \gamma_\mu L u_\alpha) + \cdots + (\overline{b_\beta} \gamma_\mu L b_\alpha)],$$

$$O_5 = (\overline{s_\alpha} \gamma^\mu L b_\alpha)[(\overline{u_\beta} \gamma_\mu R u_\beta) + \cdots + (\overline{b_\beta} \gamma_\mu R b_\beta)],$$

$$O_6 = (\overline{s_\alpha} \gamma^\mu L b_\beta)[(\overline{u_\beta} \gamma_\mu R u_\alpha) + \cdots + (\overline{b_\beta} \gamma_\mu R b_\alpha)],$$

$$O_7 = \frac{e}{16\pi^2} m_b (\overline{s_\alpha} \sigma^{\mu\nu} R b_\alpha) F_{\mu\nu},$$

$$O_8 = \frac{e^2}{16\pi^2} (\overline{s_\alpha} \gamma^\mu L b_\alpha) \overline{l} \gamma_\mu l,$$

$$O_9 = \frac{e^2}{16\pi^2} (\overline{s_\alpha} \gamma^\mu L b_\alpha) \overline{l} \gamma_\mu \gamma_5 l,$$

$$O_{10} = \frac{g}{16\pi^2} m_b (\overline{s_\alpha} \sigma^{\mu\nu} T^a_{\alpha\beta} R b_\beta) G^a_{\mu\nu}, \qquad (2)$$

where $R(L) = (1 \pm \gamma_5)/2$. Here $F_{\mu\nu}$ and $G^a_{\mu\nu}$ are the electromagnetic and strong interaction field strength tensors, and e and g are the corresponding coupling constants, respectively. In the standard model with the dimension six operator basis, O_1 and O_2 are current-current operators, $O_3, \ldots, O_6$ are usually QCD penguin operators, O_7 and O_{10} are magnetic penguin operators, and O_8 and O_9 are semileptonic electroweak operators. The C_i in Eq. (1) are the corresponding Wilson coefficients. They are calculated first at a renormalization scale M_W and then scaled down to a mass of order m_b using the renormalization group equations. It is known that the coefficients $C_3, \ldots, C_6$ are small and, thus, the contributions of the corresponding operators can be neglected. Moreover, the operator O_{10} does not contribute to the decay of $b \to s l^+ l^-$. Hence, the relevant operators of our study are O_1, O_2, O_7, O_8, and O_9, respectively. At the scale $\mu = M_W$, one has [9]

$$
\begin{aligned}
C_1(M_W) &= 0, \\
C_2(M_W) &= -1, \\
C_7(M_W) &= \tfrac{1}{2} A(x), \\
C_8(M_W) &= \frac{1}{\sin^2\theta_W} B(x) + \frac{-1+4\sin^2\theta_W}{\sin^2\theta_W} C(x) \\
&\quad + D(x) - \frac{4}{9}, \\
C_9(M_W) &= \frac{-1}{\sin^2\theta_W} B(x) + \frac{1}{\sin^2\theta_W} C(x),
\end{aligned} \tag{3}
$$

where $x = m_t^2/M_W^2$ and

$$
\begin{aligned}
A(x) &= x\left[\frac{2x^2/3+5x/12-7/12}{(x-1)^3} - \frac{3x^2/2-x}{(x-1)^4}\ln x\right], \\
B(x) &= \frac{1}{4}\left[\frac{-x}{x-1} + \frac{x}{(x-1)^2}\ln x\right], \\
C(x) &= \frac{x}{4}\left[\frac{x/2-3}{x-1} + \frac{3x/2+1}{(x-1)^2}\ln x\right], \\
D(x) &= \frac{-19x^3/36+25x^2/36}{(x-1)^3} \\
&\quad + \frac{-x^4/6+5x^3/3-3x^2+16x/9-4/9}{(x-1)^4}\ln x.
\end{aligned} \tag{4}
$$

Because O_1 and O_2 produce dilepton via virtual (vector) photon, they can be incorporated into O_8 in the later calculation. From Eq. (1) and Eq. (2), we obtain the amplitude for the inclusive process $B \to X_s l^+ l^-$ [5]:

$$
\begin{aligned}
M = \frac{G_F \alpha}{\sqrt{2}\pi} \lambda_t (C_8^{\text{eff}} \overline{s_L} \gamma_\mu b_L \overline{l} \gamma^\mu l + C_9 \overline{s_L} \gamma_\mu b_L \overline{l} \gamma^\mu \gamma_5 l \\
-2 C_7 m_b \overline{s_L} i \sigma_{\mu\nu} \frac{q^\nu}{q^2} b_R \overline{l} \gamma^\mu l).
\end{aligned} \tag{5}
$$

The coefficients C_8^{eff}, C_9, and C_7 are given by [10]

$$
\begin{aligned}
C_7(m_b) &= \eta^{-16/23}[C_7(M_W) - \tfrac{58}{135}(\eta^{10/23}-1)C_2(M_W) \\
&\quad - \tfrac{29}{189}(\eta^{28/23}-1)C_2(M_W)], \\
C_8^{\text{eff}}(m_b) &= C_8(m_b) + [3C_1(m_b) + C_2(m_b)] \\
&\quad \times \left(h(\hat{m}_c, \hat{s}) - \frac{3}{\alpha^2}\kappa \right. \\
&\quad \left. \times \sum_{V_i = J/\psi, \psi'} \frac{\pi \Gamma(V_i \to ll) M_{V_i}}{q^2 - M_{V_i} + i M_{V_i} \Gamma_{V_i}} \right), \\
C_9(m_b) &= C_9(M_W),
\end{aligned} \tag{6}
$$

where

$$
\begin{aligned}
C_1(m_b) &= \tfrac{1}{2}(\eta^{-6/23} - \eta^{12/23}) C_2(M_W), \\
C_2(m_b) &= \tfrac{1}{2}(\eta^{-6/23} + \eta^{12/23}) C_2(M_W), \\
C_8(m_b) &= C_8(M_W) + \frac{4\pi}{\alpha_s(M_W)} \left[-\frac{4}{33}[1 - \eta^{-11/23}) \right. \\
&\quad \left. + \frac{8}{87}(1 - \eta^{-29/23}) \right] C_2(M_W).
\end{aligned} \tag{7}
$$

Here $\hat{m}_c = m_c/m_b$, $\eta = \alpha_s(m_b)/\alpha_s(M_W)$, $\alpha = e^2/4\pi$, $\hat{s} = q^2/m_b^2$ with q^2 being the invariant mass of the dilepton, and $h(\hat{m}_c, \hat{s})$ arising from the one-loop contributions of O_1 and O_2 is given by

$$
\begin{aligned}
h(z, \hat{s}) &= -\tfrac{4}{9} \ln z^2 + \tfrac{8}{27} + \tfrac{4}{9} y - \tfrac{2}{9}(2+y)\sqrt{|1-y|} \\
&\quad \times \left[\Theta(1-y)\left(\ln \frac{1+\sqrt{1-y}}{1-\sqrt{1-y}} + i\pi \right) \right. \\
&\quad \left. + \Theta(y-1) 2 \arctan\left(\frac{1}{\sqrt{y-1}} \right) \right],
\end{aligned} \tag{8}
$$

where $y \equiv 4z^2/\hat{s}$. The last term of $C_8^{\text{eff}}(m_b)$ in Eq. (6) is the long-distance (LD) contribution mainly due to the J/ψ and ψ' resonances [11], and the factor κ must be chosen such that $\kappa[3C_1(m_b) + C_2(m_b)] \simeq -1$ in order to reproduce correctly the branching ratio [4]

$$
\mathcal{B}(B \to J/\psi X \to X l\overline{l}) = \mathcal{B}(B \to J/\psi X) \mathcal{B}(J/\psi \to l\overline{l}). \tag{9}
$$

III. LONGITUDINAL LEPTON POLARIZATION IN $B \to K l^+ l^-$

The hadronic matrix elements of the operators O_1, O_2, O_7, O_8, and O_9 between the B meson and the pseudoscalar K meson are given in terms of form factors as [6,7,10,12]

$$
\langle p_K | \overline{s} \gamma_\mu (1 \mp \gamma_5) b | p_B \rangle = F_+(q^2) P_\mu + F_-(q^2) q_\mu, \tag{10}
$$

$$
\begin{aligned}
&\langle p_K | \overline{s i} \sigma_{\mu\nu} q^\nu (1 \pm \gamma_5) b | p_B \rangle \\
&= \frac{1}{m_B + m_K} [P_\mu q^2 - (m_B^2 - m_K^2) q_\mu] F_T(q^2),
\end{aligned} \tag{11}
$$

where $P \equiv p_K + p_B$ and $q \equiv p_K - p_B$. In this paper, we use the form factors given by the relativistic constituent quark model in Refs. [6] and [7]. The model is based on the light front formalism, in which the form factors F_+, F_-, and F_T are taken to be approximately as

$$F(q^2) = \frac{F(0)}{1 - q^2/\Lambda_1^2 + q^4/\Lambda_2^4}. \tag{12}$$

Here the parameters Λ_1 and Λ_2 are determined by the first and second derivative of $F(q^2)$ at $q^2=0$. The values of $F(0)$ and Λ_1, Λ_2 for the various form factors used can be found in Table I of Ref. [7]. When lepton mass (for $l=e$ or μ) is negligible, q_μ terms in Eqs. (10) and (11) give no contributions to the decay rate and which is the case studied in Refs. [6] and [7]. However, the contributions have to be considered for the τ dileptonic channel of $B \to K\tau^+\tau^-$. In this case, one cannot take F_- as zero. The form factor F_- has been examined in Ref. [13], and it is found to be

$$F_- \simeq -F_+(m_B^2 + m_K^2)/(m_B^2 - m_K^2). \tag{13}$$

F_- can be extracted from other methods, such as that from heavy quark symmetry (HQS) [14]. In this method, one has

$$F_- = F_+ + \frac{2 m_B F_T}{m_B + m_K}. \tag{14}$$

The differential decay rate is given by

$$\frac{d\Gamma(B\to K l^+ l^-)}{d\hat{s}} = \frac{G_F^2 |\lambda_t|^2 m_B^5 \alpha^2}{3\cdot 2^9 \pi^5} \mathcal{D}\phi^{1/2}\left[\left[\left|C_8^{\rm eff}F_+ - \frac{2C_7F_T}{1+\sqrt{r}}\right|^2 + |C_9F_+|^2\right]\phi\left(1+\frac{2t}{\hat{s}}\right) + 12|C_9|^2 t\left[(1+r-\hat{s}/2)F_+^2 + (1-r)F_+F_- + \frac{1}{2}\hat{s}F_-^2\right]\right], \tag{15}$$

where

$$t = m_l^2/m_B^2, \quad r = m_K^2/m_B^2, \quad \hat{s} = q^2/m_B^2,$$

$$\phi = (1-r)^2 - 2\hat{s}(1+r) + \hat{s}^2, \quad \mathcal{D} = \sqrt{1-4t/\hat{s}}. \tag{16}$$

We note that the rate in Eq. (15) is a general form with $m_l \neq 0$ and it agrees with that given in Ref. [15]. The differential branching ratios, $dB/d\hat{s} \equiv d\Gamma/(\Gamma_{\rm tot} d\hat{s})$, as a function of $\hat{s}$ for $B \to K\mu^+\mu^-$ and $B \to K\tau^+\tau^-$ decays are displayed in Figs. 1(a) and 1(b), respectively. Here we have used that $m_t = 180$ GeV and $\Gamma_{\rm tot} \simeq m_b^5 G_F^2 |V_{cb}|^2/(64\pi^3)$ with $m_B \simeq m_b = 5$ GeV. The dashed and solid lines in Figs. 1(a) and 1(b) represent the results with and without the LD contributions from the resonance states, respectively. It is noted that the solid line for the decay of $B \to K\tau^+\tau^-$ in Fig. 1(b) is similar to the corresponding figure in Ref. [15].

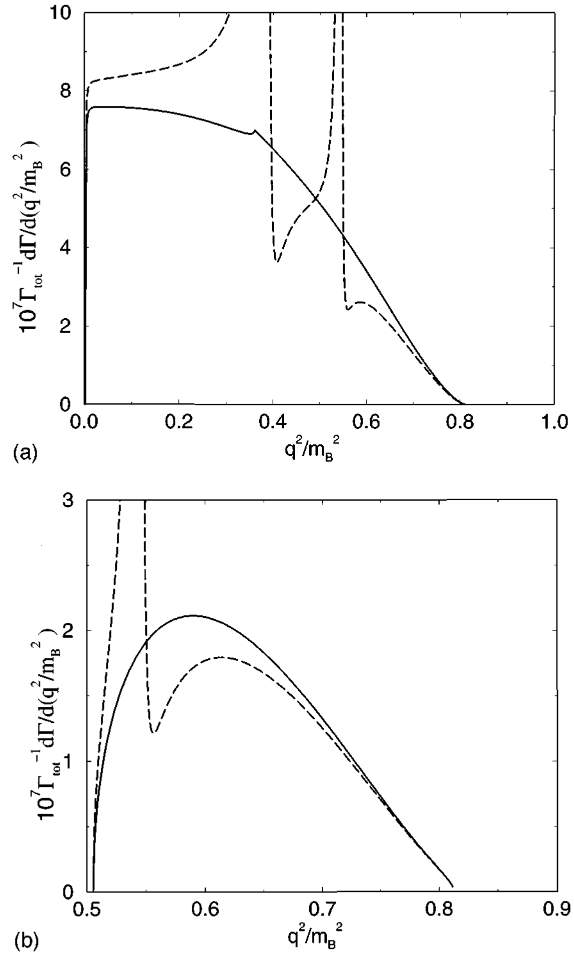

FIG. 1. Differential branching ratios for (a) $B \to K\mu^+\mu^-$ and (b) $B \to K\tau^+\tau^-$ as a function of $\hat{s} = q^2/m_B^2$ with $m_t = 180$ GeV. The dashed and solid lines correspond to the results with and without the LD contributions from the resonance states, respectively.

For the dilepton decays of the B meson, the longitudinal polarization asymmetry (LPA) of the lepton, P_L, is defined as

$$P_L(\hat{s}) = \frac{d\Gamma_{h=-1}/d\hat{s} - d\Gamma_{h=1}/d\hat{s}}{d\Gamma_{h=-1}/d\hat{s} + d\Gamma_{h=1}/d\hat{s}}, \tag{17}$$

where $h=+1$ (-1) means right (left) handed l^- in the final state and $d\Gamma/d\hat{s}$ means the differential decay rate of the B meson. In the standard model, the polarization asymmetries in Eq. (17) for $B \to K l^+ l^-$ and $B \to K^* l^+ l^-$ come from the interference of the vector or magnetic moment and axial-vector operators. For the decay of $B \to K l^+ l^-$, it is found that

$$P_L(\hat{s}) = 2\mathcal{D}\phi F_+ C_9\left(F_+ {\rm Re}C_8^{\rm eff} - 2\frac{F_T}{1+\sqrt{r}}C_7\right) \Bigg/ R, \tag{18}$$

where R is given by

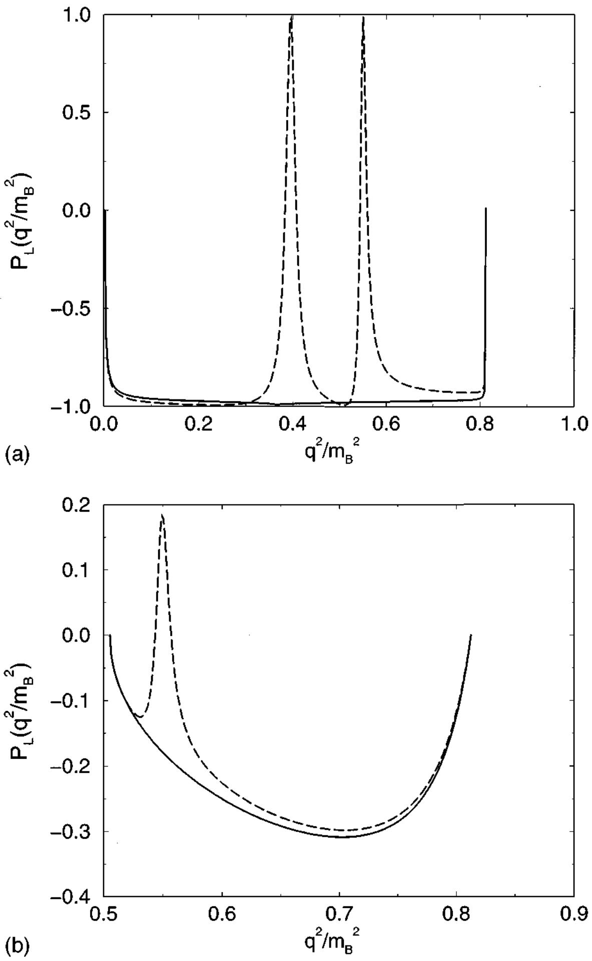

FIG. 2. Longitudinal polarization asymmetries of (a) the muon in $B\to K\mu^+\mu^-$ and (b) the τ in $B\to K\tau^+\tau^-$ as a function of $\hat{s}=q^2/m_B^2$ with $m_t=180$ GeV. Legend is the same as in Fig. 1.

$$R=\left(\left|C_8^{\rm eff}F_+-\frac{2C_7F_T}{1+\sqrt{r}}\right|^2+|C_9F_+|^2\right)\phi\left(1+\frac{2t}{\hat{s}}\right)+12|C_9|^2t[(1+r-\hat{s}/2)F_+^2+(1-r)F_+F_-+\tfrac{1}{2}\hat{s}F_-^2]. \tag{19}$$

In Figs. 2(a) and 2(b), taking $m_t=180$ GeV, we show the LPAs for $B\to Kl^+l^-$ as a function of $\hat{s}$ with $l=\mu$ and τ, respectively. The LPAs for both cases vanish at the thresholds and oscillate in the resonance regions of $q^2\simeq M^2_{\psi(\psi')}$ and they also reach zero at the end points. However, we note that if $m_l=0$, the LPA in Eq. (18) has a finite value at the maximal end point of $\hat{s}_{\rm max}=(1-m_K/m_B)^2$. From Figs. 2(a) and 2(b), we find that, away from the resonance regions and end points, $P_L(\hat{s})$ of the muon dilepton channel is around -0.9, while that of the τ is between -0.2 and -0.3 for $0.6\leqslant\hat{s}\leqslant0.8$. The average values of the LPAs are -0.8 and -0.2, respectively. Finally, we remark that Figs. 1 and 2 were drawn by using the form factor F_- in Eq. (13). The curves in Figs. 1 and 2 are nearly unchanged if one uses the form factor of F_- formulated in Eq. (14).

IV. LONGITUDINAL LEPTON POLARIZATION IN $B\to K^*l^+l^-$

The hadronic matrix elements of the operators in Eq. (5) between the external states B and K^* for the exclusive decay of $B\to K^*l^+l^-$ are [7]

$$\begin{aligned}\langle p_{K^*}|\bar{s}\gamma_\mu(1\mp\gamma_5)b|p_B\rangle=&\frac{1}{m_B+m_{K^*}}[iV(q^2)\epsilon_{\mu\nu\alpha\beta}\epsilon^{*\nu}P^\alpha q^\beta\\&\pm A_0(q^2)(m_B^2-m_{K^*}^2)\epsilon^*_\mu\pm A_+(q^2)\\&\times(\epsilon^*P)P_\mu\pm A_-(q^2)(\epsilon^*P)q_\mu],\end{aligned} \tag{20}$$

$$\begin{aligned}\langle p_{K^*}|\bar{s}i\sigma_{\mu\nu}q^\nu(1\pm\gamma_5)b|p_B\rangle=&ig(q^2)\epsilon_{\mu\nu\alpha\beta}\epsilon^{*\nu}P^\alpha q^\beta\\&\pm a_0(q^2)(m_B^2-m_K^2)\\&\times\left[\epsilon^*_\mu-\frac{1}{q^2}(\epsilon^*q)q_\mu\right]\\&\pm a_+(q^2)(\epsilon^*P)\\&\times\left[P_\mu-\frac{1}{q^2}(Pq)q_\mu\right],\end{aligned} \tag{21}$$

where $P=p_B+p_{K^*}$ and $q=p_B-p_{K^*}$. The terms corresponding to A_- in Eqs. (20) and (21) are important only for the mode of $B\to K^*\tau^+\tau^-$. Similar to Eq. (13), one has that [13]

$$A_-\simeq-A_+(m_B^2+m_{K^*}^2)/(m_B^2-m_{K^*}^2). \tag{22}$$

From Eqs. (5), (20), and (21) we obtain the differential decay rate of $B\to K^*l^+l^-$ as

$$\frac{d\Gamma(B\to K^*l^+l^-)}{d\hat{s}}=\frac{G_F^2m_B^5|\lambda_t|^2}{8\pi^3}\frac{\alpha^2}{16\pi^2}\mathcal{D}\phi^{1/2}\times\left[\left(1+\frac{2t}{\hat{s}}\right)\left(\frac{\hat{s}}{m_B^2}\alpha+\frac{\beta}{3}\phi\right)+t\delta\right], \tag{23}$$

where the definitions of α and β are the same as in Ref. [10] with choosing zero values of $C_{7'}$, $C_{8'}$, and $C_{9'}$ and

$$\begin{aligned}\delta=\frac{|C_9|^2}{2(1+\sqrt{r})^2}\Big\{&-2\phi|V|^2-3(1-r)^2|A_0|^2\\&+\frac{\phi}{4r}[2(1+r)-\hat{s}]|A_+|^2+\frac{\phi\hat{s}}{4r}|A_-|^2\\&+\frac{\phi}{2r}(1-r){\rm Re}(A_0A_+^*+A_0A_-^*+A_+A_-^*)\Big\}.\end{aligned} \tag{24}$$

The forms of r and ϕ are the same as that in Eq. (16) with the replacement $m_K\to m_{K^*}$. The differential branching ratios as a function of $\hat{s}$ for $B\to K^*\mu^+\mu^-$ and $B\to K^*\tau^+\tau^-$ de-

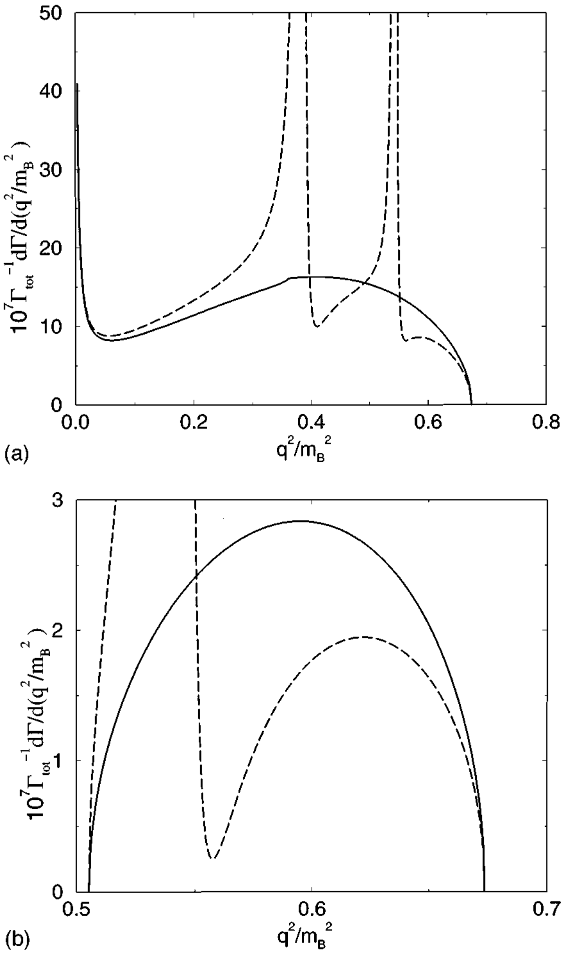

FIG. 3. Same as Fig. 1 but for (a) $B \to K^* \mu^+ \mu^-$ and (b) $B \to K^* \tau^+ \tau^-$.

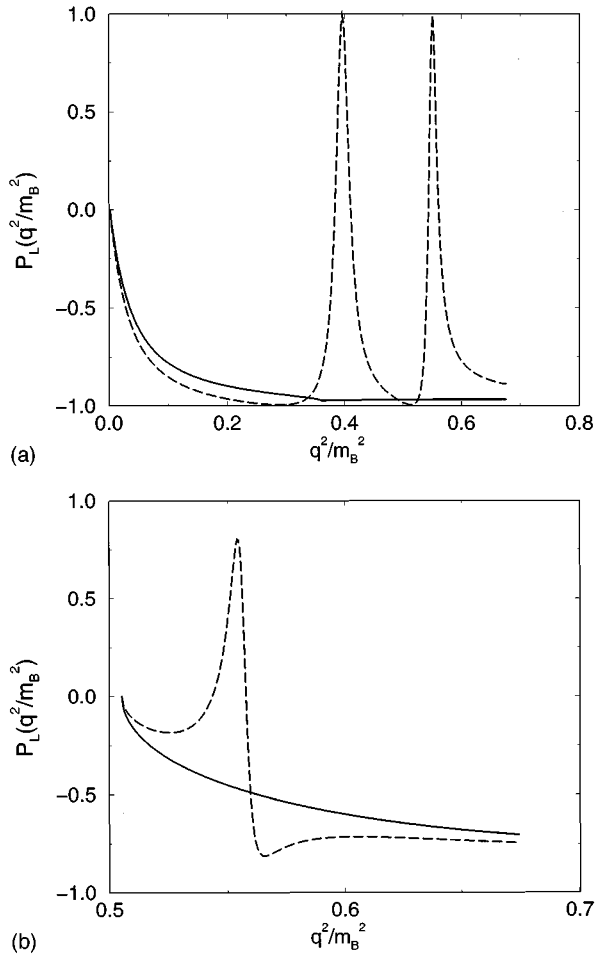

FIG. 4. Same as Fig. 2 but for (a) $B \to K^* \mu^+ \mu^-$ and (b) $B \to K^* \tau^+ \tau^-$.

cays are shown in Figs. 3(a) and 3(b), respectively, with $m_t = 180$ GeV. In Eq. (23), the lepton mass is kept to be nonzero and it is consistent with the result given in [10] for the $B \to K^* l^+ l^-$ decay by taking the limit of $m_l = 0$.

From Eq. (17), the LPA in $B \to K^* l^+ l^-$ is found to be

$$P_L(\hat{s}) = \frac{\mathcal{D}}{3} \frac{C_9}{1+\sqrt{r}} \left(\frac{\mathrm{Re}\, C_8^{\mathrm{eff}}}{1+\sqrt{r}} S_{\mathrm{av}} - \frac{2C_7}{\hat{s}} S_{\mathrm{at}} \right) \Bigg/ R, \tag{25}$$

where

$$R = \left(1 + \frac{2t}{\hat{s}}\right)\left(\frac{\hat{s}}{m_B^2}\alpha + \frac{\beta}{3}\phi\right) + t\delta,$$

$$S_{\mathrm{av}} = V^2 F_V + A_0^2 F_A^0 + A_+^2 F_A^+ + A_0 A_+ F_A^{0+},$$

$$S_{\mathrm{at}} = gVF_V + a_0 A_0 F_A^0 + \frac{a_0 A_+ + a_+ A_0}{2} F_A^{0+} + a_+ A_+ F_A^+, \tag{26}$$

with

$$F_V = \hat{s}\phi,$$

$$F_A^0 = \frac{1}{8r}(1-r)^2[(1-r)^2 - 2\hat{s} + 10r\hat{s} + \hat{s}^2],$$

$$F_A^+ = \frac{1}{8r}\phi^2,$$

$$F_A^{0+} = \frac{1}{4r}\phi(1-r)(1-r-\hat{s}). \tag{27}$$

In Figs. 4(a) and 4(b), we plot the LPAs as a function of $\hat{s}$ for $B \to K^* \mu^+ \mu^-$ and $K^* \tau^+ \tau^-$, respectively, with $m_t = 180$ GeV. From Fig. 4, we see that $P_L(B \to K^* l^+ l^-)$ for both $l = \mu$ and τ vanish at the minimal end point of $\hat{s} = 4m_l^2/m_B^2$ due to the kinematic factor $\mathcal{D}$ in Eq. (25). However, in contrast with the cases in $B \to K l^+ l^-$ shown in Figs. 2(a) and 2(b), $P_L(B \to K^* l^+ l^-)$ with $l = \mu$ and τ are not zero at the maximum of $\hat{s}_{\max} = (1 - m_{K^*}/m_B)^2$. This is due to the fact that the numerator in Eq. (18) has extra zero from ϕ at $\hat{s}_{\max}$ whereas it does not contain ϕ in Eq. (25). Similar to the cases of $B \to K l^+ l^-$, $P_L(B \to K^* \mu^+ \mu^-)$ has a large negative value over most of the allowed range of $\hat{s}$, with an average value $\langle P_L(B \to K^* \mu^+ \mu^-) \rangle = -0.7$ while for the $\tau^+ \tau^-$ channel, the average τ polarization is

$\langle P_L(B\to K^*\tau^+\tau^-)\rangle=-0.5$. Finally, we note that our results in Figs. 3 and 4 are insensitive to the values of A_- predicted in the different form factor models.

V. CONCLUSIONS

We have investigated the longitudinal polarization asymmetries of the muon and τ in the exclusive processes of $B\to Kl^+l^-$ and $B\to K^*l^+l^-$. The average values of the polarization asymmetries of the muon and τ are found to be -0.8 and -0.2 for $B\to K\mu^+\mu^-$ and $B\to K\tau^+\tau^-$, and -0.7 and -0.5 for $B\to K^*\mu^+\mu^-$ and $B\to K^*\tau^+\tau^-$, respectively. These polarization asymmetries provide valuable information on the flavor changing loop effects in the standard model. From Figs. 1 and 3, we find that the total integrated branching ratios of $B\to K^{(*)}l^+l^-$ are 0.5 $(1.4)\times10^{-6}$ and 1.3 $(2.2)\times10^{-7}$ with $l=\mu$ and τ, respectively. Experimentally, to measure an asymmetry A of a decay with the branching ratio B at the $n\sigma$ level, the required number of events is $N=n^2/(BA^2)$. For example, to observe the τ polarizations at the both exclusive channels of $B\to K^{(*)}\tau^+\tau^-$, one needs at least $1.8\times10^7 n^2$ $B\bar{B}$ decays. Therefore, at the B factories under construction, some of the asymmetries could be accessible.

ACKNOWLEDGMENTS

We thank C. W. Hwang for showing their recent results [13] on the hadronic form factors to us. This work was supported in part by the National Science Council of the Republic of China under Grant No. NSC-85-2112-M-007-031.

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