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

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We study the transverse lepton polarization (P_T) in the exclusive processes of $B \to K \ell^+ \ell^-$ and $B \to K^* \ell^+ \ell^-$. We find that the polarization is zero for $B \to K \ell^+ \ell^-$ but nonzero for $B \to K^* \ell^+ \ell^-$. The average values of $P_T(B \to K^* \ell^+ \ell^-)$ are 8.4×10^{-4} and 2.0×10^{-2} for $\ell = \mu$ and $\ell = \tau$, respectively. The polarization asymmetry is a T-odd observable, which connects with the non-Hermiticity of the effective Hamiltonian, arising mainly from the $c\bar{c}$ intermediate states. We also include the other two orthogonal components, longitudinal and normal polarizations, in our discussions. [S0556-2821(98)02107-9]

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I. INTRODUCTION

It is well known that, in the standard model, dilepton B decays such as $B \to X \ell^+ \ell^-$ do not occur at the tree level and appear only at the quantum loop level. The observations of the exclusive $B \to K^* \gamma$ as well as the inclusive $B \to X_s \gamma$ modes by the CLEO Collaboration [1] have stimulated the study of rare B decays induced by the flavor-changing neutral current $b \to s$ transition [2]. Experimentally, the upper limits on the exclusive decays $B \to K^{(*)} \ell^+ \ell^-$ from CLEO and the Collider Detector at Fermilab (CDF) [3] are less than one order of magnitude above the standard model predictions, respectively. In the standard model, the short-distance (SD) contributions to the decays involving the $b \to s$ transition are dominated by loops with the top quark, while the long-distance (LD) contributions arise mainly from the J/ψ and ψ' resonances. Thus, the rare B decays are a good probe of heavy top quark physics. Furthermore, it is also important to study the decays of $B \to X_s l^+ l^-$ due to the potential test of the effective Hamiltonian [4] and QCD corrections [5].

It is pointed out by Ali *et al.* [6] that, in the standard model, the measurement of the forward-backward asymmetry of the dileptons in the inclusive $b \to s \ell^+ \ell^-$ decay provides information on the short-distance contribution dominated by the top quark loops. Other interesting asymmetries, such as the transverse, normal and longitudinal lepton polarization asymmetries in the inclusive processes $b \to s \ell^+ \ell^-$, denoted by P_T, P_N, and P_L, respectively, are also important observables. These asymmetries have been recently studied by Krüger *et al.* [7] and Hewett [8]. The P_T is an T-odd observable, arising mainly from the non-Hermiticity of the effective Hamiltonian. Since P_N and P_T are both proportional to the ratio of m_l/m_B, the $\tau^+ \tau^-$ channel is especially noticeable.

On the other hand, for the exclusive processes of $B \to K^{(*)} \ell^+ \ell^-$, the longitudinal polarization asymmetry has been studied by Ref. [9]. In this report, we will investigate the remaining two polarization asymmetries P_T and P_N. We will show that the forward-backward and transverse lepton polarization asymmetries of the exclusive decays $B \to M l^+ l^-$ are identically zero where M is a pseudoscalar meson such as π and K, but nonzero when M is a vector meson such as ρ and K^*. But the longitudinal and normal lepton polarizations for the exclusive modes are nonzero for both cases of the pseudoscalar and vector mesons. In our discussions, we will use the results of the relativistic quark model by the light-front formalism [10,11] on the form factors in the hadronic matrix elements between the B 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. The form factors between hadrons are also discussed. We study the polarization asymmetries for the exclusive decays of $B \to K^{(*)} \ell^+ \ell^-$ in Sec. III. Our conclusions are summarized in Sec. IV.

II. EFFECTIVE HAMILTONIAN AND FORM FACTORS

The effective Hamiltonian relevant to $b \to s l^+ l^-$ based on an operator product expansion is given by [12]

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

where G_F denotes the Fermi constant, $\lambda_t = V_{tb} V_{ts}^*$ is the product of the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements. The operators $O_i(\mu)$ and the analytic expressions for the Wilson coefficients are given in Ref. [12]. From Eq. (1) , we obtain the amplitude for the inclusive process $B \to X_s \ell^+ \ell^-$ [6,8,12,13] as

$$M = \frac{G_F \alpha}{\sqrt{2} \pi} \lambda_t \bigg(C_8^{eff} \bar{s}_L \gamma_\mu b_L \bar{l} \gamma^\mu l + C_9 \bar{s}_L \gamma_\mu b_L \bar{l} \gamma^\mu \gamma_5 l - 2 C_7 m_b \bar{s}_L i \sigma_{\mu\nu} \frac{q^\nu}{q^2} b_R \bar{l} \gamma^\mu l \bigg). \tag{2}$$

Among the Wilson coefficients in Eq. (2), the most important one is C_8^{eff}, given by

$$C_8^{eff}(m_b)=C_8(m_b)+[3C_1(m_b)+C_2(m_b)]\times\left(h(\hat{m}_c,\hat{s})-\frac{3}{\alpha^2}\kappa\times\sum_{V_i=J/\psi,\psi'}\frac{\pi\Gamma(V_i\to ll)M_{V_i}}{q^2-M_{V_i}^2+iM_{V_i}\Gamma_{V_i}}\right), \quad (3)$$

with

$$C_1(m_b)=\frac{1}{2}(\eta^{-6/23}-\eta^{12/23})C_2(M_W),$$

$$C_2(m_b)=\frac{1}{2}(\eta^{-6/23}+\eta^{12/23})C_2(M_W),$$

$$C_8(m_b)=C_8(M_W)+\frac{4\pi}{\alpha_s(M_W)}\left\{-\frac{4}{33}[1-\eta^{-11/23}]+\frac{8}{87}[1-\eta^{-29/23}]\right\}C_2(M_W), \quad (4)$$

$$h(z,\hat{s})=-\frac{4}{9}\ln z^2+\frac{8}{27}+\frac{4}{9}y-\frac{2}{9}(2+y)\sqrt{|1-y|}\times\left\{\Theta(1-y)\left[\ln\left|\frac{1+\sqrt{1-y}}{1-\sqrt{1-y}}\right|-i\pi\right]+\Theta(y-1)2\arctan\left(\frac{1}{\sqrt{y-1}}\right)\right\}, \quad (5)$$

where $\hat{m}_c=m_c/m_b$, $\eta=\alpha_s(m_b)/\alpha_s(M_W)$, $\alpha=e^2/4\pi$, and $\hat{s}=q^2/m_b^2$ with q^2 being the invariant mass of the dilepton. In Eq. (5), it is found that $h(\hat{m}_c,\hat{s})$, arising from the one-loop contributions of O_1 and O_2, possesses an imaginary part $i\pi$ which appears from the associated penguin diagram. Letting H_{peng} be the effective Hamiltonian derived from the penguin diagram, it contains the form

$$H_{peng}\sim(\bar{s}_L\gamma_\mu b_L\bar{l}\gamma^\mu l)\left(F_1(\hat{m}_c^2)+F_2(\hat{m}_c^2)\times\int_0^1dx\,\ln[q^2x(1-x)-m_c^2]\right), \quad (6)$$

with $F_1(x)$ and $F_2(x)$ being analytic function of x. Clearly, it becomes complex when the invariant mass of the dilepton is above the $c\bar{c}$ threshold, i.e., $q^2>4m_c^2$ and real when it is below the threshold. The last term of $C_8^{eff}(m_b)$ in Eq. (3) is the LD contribution mainly due to the J/ψ and ψ' resonances [14], and the factor κ must be chosen such that $\kappa(3C_1(m_b)+C_2(m_b))\simeq-1$ in order to correctly reproduce the branching ratio of [6]

$$BR(B\to J/\psi X\to Xl\bar{l})=BR(B\to J/\psi X)BR(J/\psi\to l\bar{l}). \quad (7)$$

We will see later that the imaginary part of C_8^{eff} is the main source that contributes to the T-odd observable of P_T. However, the last term in Eq. (3) is close to zero if q^2 is away from the resonance regions.

In order to study the exclusive processes, we define the hadronic matrix elements $\langle p_{K(*)}|\bar{s}\gamma_\mu(1\mp\gamma_5)b|p_B\rangle$ and $\langle p_{K(*)}|\bar{s}i\sigma_{\mu\nu}q^\nu(1\pm\gamma_5)b|p_B\rangle$ as follows:

$$\langle p_K|\bar{s}\gamma_\mu(1\mp\gamma_5)b|p_B\rangle=F_+(q^2)P_\mu+F_-(q^2)q_\mu, \quad (8)$$

$$\langle p_K|\bar{s}i\sigma_{\mu\nu}q^\nu(1\pm\gamma_5)b|p_B\rangle=\frac{1}{m_B+m_K}[P_\mu q^2-(m_B^2-m_K^2)q_\mu]F_T(q^2), \quad (9)$$

$$\langle p_{K*}|\bar{s}\gamma_\mu(1\mp\gamma_5)b|p_B\rangle=\frac{1}{m_B+m_{K*}}[iV(q^2)\epsilon_{\mu\nu\alpha\beta}\epsilon^{*\nu}P^\alpha q^\beta \pm A_0(q^2)(m_B^2-m_{K*}^2)\epsilon_\mu^* \pm A_+(q^2)(\epsilon^*P)P_\mu \pm A_-(q^2)(\epsilon^*P)q_\mu], \quad (10)$$

$$\langle p_{K*}|\bar{s}i\sigma_{\mu\nu}q^\nu(1\pm\gamma_5)b|p_B\rangle=ig(q^2)\epsilon_{\mu\nu\alpha\beta}\epsilon^{*\nu}P^\alpha q^\beta \pm a_0(q^2)(m_B^2-m_K^2)\left[\epsilon_\mu^*-\frac{1}{q^2}(\epsilon^*q)q_\mu\right] \pm a_+(q^2)(\epsilon^*P)\left[P_\mu-\frac{1}{q^2}(Pq)q_\mu\right], \quad (11)$$

where $P\equiv p_B+p_{K(*)}$, and $q\equiv p_B-p_{K(*)}$. There are many models used to evaluate the above form factors [15]. In this paper, we use the form factors given by the relativistic constituent quark model [10]. Although the q_μ term does not contribute to the differential decay rate, it is important in the calculations of the various polarization asymmetries. For the form factors $G_-(=F_-,A_-)$, we relate them with $G_+(=F_+,A_+)$ by the relation [16]

$$G_-\simeq-G_+(m_B^2+m_{K(*)}^2)/(m_B^2-m_{K(*)}^2). \quad (12)$$

III. LEPTON POLARIZATIONS IN $B\to K\ell^+\ell^-$ AND $B\to K^*\ell^+\ell^-$

The lepton polarization vector can be decomposed into three orthogonal components: longitudinal, normal, and transverse. By writing the three orthogonal unit vectors as

$$\mathbf{e}_L=\frac{\vec{p}_-}{|\vec{p}_-|}, \quad (13)$$

$$\mathbf{e}_N=\frac{\vec{p}_-\times(\vec{p}_{K(*)}\times\vec{p}_-)}{|\vec{p}_-\times(\vec{p}_{K(*)}\times\vec{p}_-)|}, \quad (14)$$

$$\mathbf{e}_T=\frac{\vec{p}_{K(*)}\times\vec{p}_-}{|\vec{p}_{K(*)}\times\vec{p}_-|}, \quad (15)$$

with $\vec{p}_-$ and $\vec{p}_{K^{(*)}}$ being the three momenta of ℓ^- and $K^{(*)}$, we can define the longitudinal, normal, and transverse lepton polarization asymmetries by

$$P_i(\hat{s}) = \frac{d\Gamma(\mathbf{n}=\mathbf{e}_i)/d\hat{s} - d\Gamma(\mathbf{n}=-\mathbf{e}_i)/d\hat{s}}{d\Gamma(\mathbf{n}=\mathbf{e}_i)/d\hat{s} + d\Gamma(\mathbf{n}=-\mathbf{e}_i)/d\hat{s}}, \tag{16}$$

where $\mathbf{n}$ is the polarization vector of ℓ^-. The differential decay rate of $B \to K^{(*)}\ell^-\ell^+$ is given by

$$\frac{d\Gamma}{d\hat{s}} = \frac{1}{2}\left(\frac{d\Gamma}{d\hat{s}}\right)_{unpol}[1+(P_L\mathbf{e}_L + P_T\mathbf{e}_T + P_N\mathbf{e}_N)\cdot\mathbf{n}]. \tag{17}$$

As can be easily seen from Eqs. (15) and (17), the time-reversal operation changes only the sign of P_T. Therefore, a nonzero value of P_T implies an T-odd observable. P_L and P_N are related to the polarization in the decay plane of the lepton and kaon and thus T-even observables. The formula for P_L has been shown in Ref. [9]. Here we only discuss the polarization asymmetries of P_N and P_T, which are given by

$$P_N(\hat{s})(B\to K\ell^+\ell^-) = \frac{3\pi\hat{m}_l}{2\sqrt{\hat{s}}}\phi^{1/2}\left(2\,\mathrm{Re}\,C_9C_7\frac{F_T}{1+\sqrt{r}} - \mathrm{Re}\,C_8^{eff}C_9F_+\right) \times (F_+(1-r)+F_-\hat{s})/R_K, \tag{18}$$

$$P_T(\hat{s})(B\to K\ell^+\ell^-) = 0, \tag{19}$$

$$P_N(\hat{s})(B\to K^*\ell^+\ell^-) = \frac{\pi\hat{m}_l}{32\sqrt{\hat{s}}}\phi^{1/2}\Omega(\hat{s})/R_{K^*}, \tag{20}$$

$$P_T(\hat{s})(B\to K^*\ell^+\ell^-) = -\frac{\pi\hat{m}_l}{4}\phi^{1/2}\mathcal{D}\frac{1-\sqrt{r}}{1+\sqrt{r}}\sqrt{\hat{s}}\ \mathrm{Im}(C_8^{eff}C_9)VA_0/R_{K^*}, \tag{21}$$

where

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

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

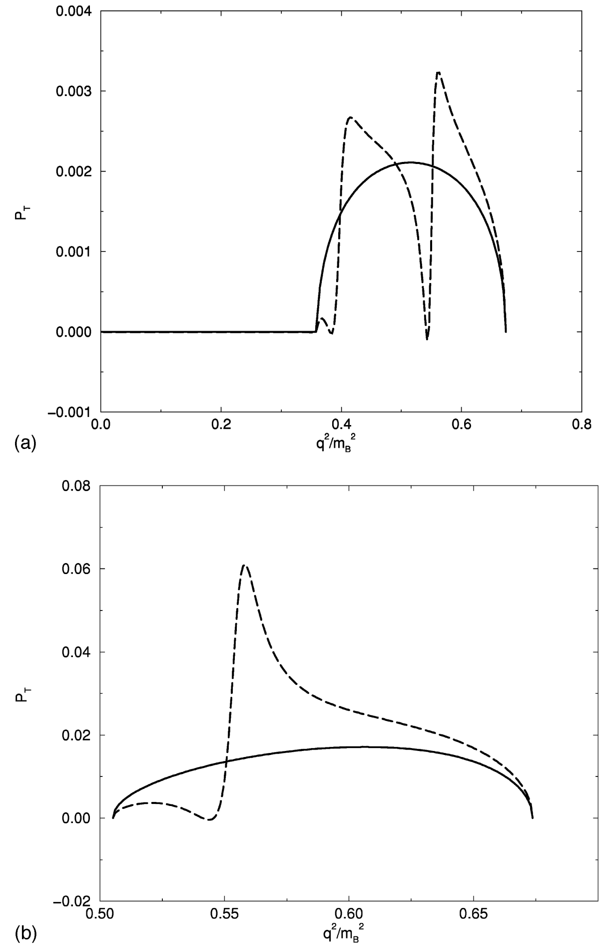

FIG. 1. Transverse polarization asymmetry of (a) the μ in $B \to K^*\mu^+\mu^-$ and (b) the τ in $B\to K^*\tau^+\tau^-$ as function of $\hat{s}$ with $m_t = 180$ GeV. The dashed and solid lines correspond to the results with and without the LD contributions from the resonance states, respectively.

$$\Omega(\hat{s}) = \frac{-8|C_8^{eff}|^2}{(1+\sqrt{r})^2}VA_0(1-r)\hat{s} - 32|C_7|^2 g a_0(1-r) + \frac{\mathrm{Re}\,C_8^{eff}C_9}{(1+\sqrt{r})^2}\frac{1}{r}[(A_+ + A_-)(-1+r) - A_-\hat{s}] \times [A_+\phi + A_0(1-r)(1-r-\hat{s})] + \frac{16\,\mathrm{Re}\,C_8^{eff}C_7}{1+\sqrt{r}} \times (Va_0 + A_0 g)(1-r) - \frac{2\,\mathrm{Re}\,C_9C_7}{1+\sqrt{r}}\frac{1}{r\hat{s}} \times [(A_+ + A_-)(-1+r) - A_-\hat{s}][a_+\phi + a_0(1-r) \times (1-r-\hat{s})], \tag{22}$$

with $r \equiv m_{K^{(*)}}^2/m_B^2$ and $t \equiv m_l^2/m_B^2$. The formulas for α, β, δ, $\mathcal{D}$, and ϕ, in Eqs. (18)–(22), can be found in Refs. [9,10], respectively. The P_T is the T-odd observable of interest. As is expected from the $\mathrm{Im}(F_V F_A^*)$ term in Eq. (3) of Ref. [17] it comes mainly from the absorptive part of C_8^{eff} induced from the loop as well as QCD corrections. Both O_1 and O_2 contribute to this absorptive part. But the latter is α_s

times smaller than the former. Because of $m_b \gg m_c$, the charm-quark loop of four-quark operator matrix element can be on shell and leads to the absorptive part. From Eq. (18), it is found that P_N for $B\to K\ell^+\ell^-$ has only two terms instead of six terms in inclusive process [7]. One of them arises from the interference between the vector and axial-vector operators, while the other from the interference between the axial-vector and magnetic moment operators. In addition, the P_T for $B\to K\ell^+\ell^-$ is zero due to the pseudoscalar-pseudoscalar transition of $B\to K$, which gives no terms proportional to $\epsilon_{\mu\nu\alpha\beta}$ unlike the case of $B\to K^*\ell^+\ell^-$, in which P_T is nontrivial since K^* is a spin-one vector meson.

The longitudinal polarization has been studied in Ref. [9]. It is found that average values of P_L of the muon and tau for $B\to K^{(*)}\mu^+\mu^-$ and $B\to K^{(*)}\tau^+\tau^-$ are $-0.8(-0.7)$ and $-0.2(-0.5)$, respectively. Similar figures can be shown for the normal polarizations in $B\to K^{(*)}\ell^+\ell^-$ ($\ell=e,\mu$) decays. We observe that the LD contributions enhance both P_L and P_N. We note that, in the neighborhood of $\hat{s}_{max}\equiv(1-m_K/m_B)^2$, $P_N(B\to K\ell^+\ell^-)$ is decreasing for $\ell=\mu$ but increasing for $\ell=\tau$. The average values of P_N of the muon and τ are found to be 0.1 and 0.6 for $B\to K\mu^+\mu^-$ and $B\to K\tau^+\tau^-$, respectively. In addition, as is expected, the contribution to P_N from the $\tau^+\tau^-$ channel is much larger than that from the $\mu^+\mu^-$ one. Unlike the case in $B\to K\ell^+\ell^-$, which contains only interference terms, there are extra square terms that contribute to P_N in $B\to K^*\ell^+\ell^-$. We note that, among P_L, P_N and P_T, P_N in $B\to K^*\ell^+\ell^-$ is the only quantity to receive a contribution from a single *vector* (O_8) or *magnetic moment* (O_7) operator. We find that $P_N(B\to K^*\ell^+\ell^-)$ for both $\ell=\mu$ and $\ell=\tau$ vanish at the maximal end point $\hat{s}_{max}\equiv(1-m_{K^*}/m_B)^2$ due to the kinematic factor ϕ in Eq. (20). The LD contribution to $P_N(B\to K^*\tau^+\tau^-)$ is always constructive relative to the SD part. However, the LD contribution to $P_N(B\to K^*\mu^+\mu^-)$ is oscillatory relative to the SD part. The average values of P_N of the muon and τ are found to be about 0.02 and 0.4 for $B\to K^*\mu^+\mu^-$ and $B\to K^*\tau^+\tau^-$, respectively. It is obvious that the $\tau^+\tau^-$ channel has more considerable contribution than the $\mu^+\mu^-$ channel.

The P_T for $B\to K^*\ell^+\ell^-$ is plotted in Figs. 1(a) and 1(b) for $l=\mu$ and τ, respectively. It is kinematically forbidden at the two end points. The LD effect in Fig. 1(b) is oscillatory relative to the SD part. Although only the interference of vector and axial-vector operators could make contributions to the $P_T(B\to K^*\ell^+\ell^-)$, it still vanishes if there is no $c\bar{c}$ resonance contribution to the coefficient C_8^{eff}. In Fig. 1(a), it can be easily found that, for $\hat{s}<4\hat{m}_c^2$, P_T is close to zero. The average values of P_T are found to be 8.4×10^{-4} and 2.0×10^{-2} for $B\to K^*\mu^+\mu^-$ and $B\to K^*\tau^+\tau^-$, respectively. We remark that our results of the T-odd transverse lepton polarization asymmetries for the exclusive B decays are consistent with those given by Ref. [7] for the inclusive ones.

IV. CONCLUSIONS

The transverse and normal polarization asymmetries for exclusive dilepton rare B decays of $B\to K\ell^+\ell^-$ and $B\to K^*\ell^+\ell^-$ are studied. We have shown that the transverse lepton polarization is zero for $B\to K\ell^+\ell^-$ but nonzero for $B\to K^*\ell^+\ell^-$. The average values of P_T are found to be 8.4×10^{-4} and 2.0×10^{-2} for $B\to K^*\mu^+\mu^-$ and $B\to K^*\tau^+\tau^-$, respectively. The transverse polarization asymmetry is a T-odd observable which connects with the non-hermiticity of the effective Hamiltonian, arising mainly from the $c\bar{c}$ intermediate states. The average values of P_N of the muon and τ are found to be 0.1 and 0.6 for $B\to K\mu^+\mu^-$ and $B\to K\tau^+\tau^-$, and 0.02 and 0.4 for $B\to K^*\mu^+\mu^-$ and $B\to K^*\tau^+\tau^-$, respectively. Finally, we remark that some of the polarization asymmetries for τ could be accessible in the future B factories. For example, experimentally, to observe the tau polarizations of (P_L,P_N,P_T) in $B\to K^*\tau^+\tau^-$ at the $n\sigma$ level, one needs at least $(1.8,2.8,10^3)\times10^7 n^2$ $B\bar{B}$ decays, respectively.

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[1] CLEO Collaboration, R. Ammar *et al.*, Phys. Rev. Lett. **71**, 674 (1993); CLEO Collaboration, M. S. Alam *et al.*, *ibid.* **74**, 2885 (1995).

[2] For a recent review, see J. L. Hewett, in *Spin Structure in High Energy Processes*, Proceedings of the SLAC Summer Institute on Particle Physics, Stanford, CA, 1993, edited by L. DePorcel and C. Dunwoodil (SLAC Report No. 444, Stanford, 1994), pp. 463–475, and references therein.

[3] CLEO Collaboration, R. Balest *et al.*, CLEO-CONF-94-4 and CDF Collaboration, C. Anway-Wiese *et al.*, Fermilab-Conf-95/201-E.

[4] T. Inami and C. S. Lim, Prog. Theor. Phys. **65**, 297 (1981).

[5] A. J. Buras and M. Münz, Phys. Rev. D **52**, 186 (1995).

[6] A. Ali, T. Mannel, and T. Morozumi, Phys. Lett. B **273**, 505 (1991).

[7] F. Kruger and L. M. Sehgal, Phys. Lett. B **380**, 199 (1996).

[8] J. L. Hewett, Phys. Rev. D **53**, 4964 (1996).

[9] C. Q. Geng and C. P. Kao, Phys. Rev. D **54**, 5636 (1996).

[10] W. Jaus and D. Wyler, Phys. Rev. D **41**, 3405 (1990).

[11] W. Jaus, Phys. Rev. D **41**, 3394 (1990).

[12] B. Grinstein, M. B. Wise, and M. Savage, Nucl. Phys. **B319**, 271 (1989).

[13] M. Misiak, Nucl. Phys. **B393**, 23 (1993); **B439**, 461(E) (1995).

[14] 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. D'Donnell and H. K. K. Tung, *ibid.* **43**, R2067 (1991).

[15] N. G. Deshpande and J. Trampetic, Phys. Rev. Lett. **60**, 2583 (1988); G. Burdman, Phys. Rev. D **52**, 6400 (1995); P. Colangelo *et al.*, *ibid.* **53**, 3672 (1996); W. Roberts, *ibid.* **54**, 863 (1996).

[16] H.-Y. Cheng, C.-Y. Cheung, and C.-W. Hwang, Phys. Rev. D **55**, 1559 (1997) and (private communication); C. Y. Cheung, C. W. Hwang, and W. M. Zhang, Z. Phys. C **75**, 657 (1997).

[17] P. Agrawal, J. N. Ng, G. Bélanger, and C. Q. Geng, Phys. Rev. D **45**, 2383 (1992).