Exclusive semileptonic rare decays $B\rightarrow(K,K^*)l^+l^-$ in supersymmetric theories

Qi-Shu Yan*
Physics Department of Tsinghua University, People's Republic of China

Chao-Shang Huang,† Wei Liao,‡ and Shou-Hua Zhu§
Institute of Theoretical Physics, the Chinese Academy Science, People's Republic of China

(Received 1 May 2000; published 11 October 2000)

The invariant mass spectrum, forward-backward asymmetry, and lepton polarizations of the exclusive processes $B\rightarrow K(K^*)l^+l^-, l=\mu,\tau$ are analyzed in a supersymmetric context. Special attention is paid to the effects of neutral Higgs bosons (NHB's). Our analysis shows that the branching ratio of the process $B\rightarrow K\mu^+\mu^-$ can be quite largely modified by the effects of neutral Higgs bosons and the forward-backward asymmetry would not vanish. For the process $B\rightarrow K^*\mu^+\mu^-$, the lepton transverse polarization is quite sensitive to the effects of NHB's, while the invariant mass spectrum, forward-backward asymmetry, and lepton longitudinal polarization are not. For both $B\rightarrow K\tau^+\tau^-$ and $B\rightarrow K^*\tau^+\tau^-$, the effects of NHB's are quite significant. The partial decay widths of these processes are also analyzed, and our analysis manifests that, even taking into account the theoretical uncertainties in calculating weak form factors, the effects of NHB's could make supersymmetry show up.

PACS number(s): 13.20.He, 12.60.Jv, 13.25.Hw

I. INTRODUCTION

The inclusive rare processes $b\rightarrow X_s l^+l^-, l=e, \mu, \tau$ have been intensively studied in the literature [1–12]. As one kind of the flavor changing neutral current processes, they are sensitive to the fine structure of the standard model and to possible new physics as well, and are expected to shed light on the existence of new physics before the possible new particles are produced at colliders.

It is well known that the invariant mass spectrum, forward-backward asymmetries (FBA's), and lepton polarizations are important observables to probe new physics, while the first two observables are mostly analyzed. About lepton polarizations, it is known that, due to the smallness of the mass of it, electron polarizations are very difficult to measure experimentally. So only the lepton polarizations of the muon and tau are considered in the literature [10,12–14]. The longitudinal polarization of tau in $B\rightarrow X_s\tau^+\tau^-$ has been calculated in the standard model (SM) and several new physics scenarios [10]. For $B\rightarrow X_s l^+l^-$ $(l=\mu,\tau)$, the polarization of the lepton in the SM is analyzed in [12] and it is pointed out that, for the μ channel, the only significant component is the longitudinal polarization (P_L), while all three components are sizable in the τ channel. The analysis has been extended to supersymmetric models (SUSY) and a CP softly broken two-Higgs-doublet model in Refs. [13] and [14], respectively. Reference [5] also gives a general model-independent analysis of the lepton polarization asymmetries in the process $B\rightarrow X_s\tau^+\tau^-$ and it is found that the contribution from $C_{LRLR}+C_{LRRL}$ is much larger than other scalar-type interactions.

Compared with the inclusive processes $B\rightarrow X_s l^+l^-, l=e, \mu, \tau$, the theoretical study of the exclusive processes $B\rightarrow K(K^*)l^+l^-$ is relatively hard. For inclusive semileptonic decays of B, the decay rates can be calculated in heavy quark effective theory (HQET) [15]. However, for exclusive semileptonic decays of B, to make theoretical predictions, additional knowledge of decay form factors is needed, which is related with the calculation of hadronic transition matrix elements. Hadronic transition matrix elements depend on the nonperturbative properties of QCD, and can only be reliably calculated by using a nonperturbative method. The form factors for B decay into $K^{(*)}$ have been computed with different methods such as quark models [16], Shifman-Vainshtain-Zakharov (SVZ) QCD sum rules [17], light cone sum rules (LCSR's) [18–22]. Compared to the lattice approach which mainly deal with the form factors at small recoil, the QCD sum rules on the light cone can complementarily provide information on the form factors at smaller values of $\hat{s}$. And they are consistent with perturbative QCD and the heavy quark limit. In this work, we will use the weak decay form factors calculated by using the technique of the light cone QCD sum rules and given in [23].

A upper limit on the branching ratio of $B^0\rightarrow K^{0*}\mu^+\mu^-$ has been recently given by CLEO [24]:

$$BR(B^o\rightarrow K^{0*}\mu^+\mu^-)<4.0\times10^{-6}, \qquad (1.1)$$

and they will be precisely measured at B factories. These exclusive processes are quite worthy of intensive study and have attracted much attention [23,25–36]. In Ref. [23], by using improved theoretical calculations of the decay form factors in the light cone QCD sum rule approach, dilepton invariant mass spectra and the FBA's of these exclusive decays are analyzed in the standard model and a number of popular variants of the supersymmetric models. However, as the author claimed, the effects of neutral Higgs exchanges are neglected. For exclusive processes, as pointed out in

*Email address: yanqs@itp.ac.cn; qsyan@mail.tsinghua.edu.cn
†Email address: csh@itp.ac.cn
‡Email address: liaow@itp.ac.cn
§Email address: huald@itp.ac.cn

[36], the polarization asymmetries of μ and τ for $B \to K^*\mu^+\mu^-$ and $B\to K^*\tau^+\tau^-$ are also accessible at the B factories under construction. In Ref. [34], the lepton polarizations and CP-violating effects in $B\to K^*\tau^+\tau^-$ are analyzed in SM and two-Higgs-doublet models.

As pointed in Refs. [3,4], in two-Higgs-doublet models (2HDM) and SUSY models, the neutral Higgs boson could contribute largely to the inclusive processes $B\to X_s l^+l^-, l=\mu, \tau$ and greatly modify the branching ratio and FBA in the large $\tan\beta$ case. The effects of the neutral Higgs boson in the 2HDM to polarizations of τ in $B\to K\tau^+\tau^-$ are analyzed in [35], and it was found that polarizations of the charged final lepton are very sensitive to the $\tan\beta$.

In this paper, we will investigate the exclusive decay $B \to K(K^*)l^+l^-, l=\mu, \tau$ in SUSY models. We shall evaluate branching ratios and FBA's with an emphasis on the effects of the neutral Higgs boson and analyze lepton polarizations in the minimal supersymmetric standard model (MSSM). According to the analysis of [29], different sources of the vector current could manifest themselves in different regions of phase space, for the very low $\hat{s}$ the photonic penguin diagram dominates, while the Z penguin diagram and W box diagram becomes important towards high $\hat{s}$. In order to search the regions of $\hat{s}$ where neutral Higgs bosons could greatly contribute, we analyze the partial decay widths of these two processes. Beside that they are accessible to B factories, our motivation also is based on the fact that to the inclusive processes $B\to X_s l^+l^-, l=\mu, \tau$, neutral Higgs boson could make quite a large contribution at certain large $\tan\beta$ regions of parameter space in SUSY models, since part of supersymmetric contributions is proportional to $\tan^3\beta$ [4]. Such regions considerably exist in supergravity (SUGRA) and M-theory inspired models [37]. We also analyze the effects of neutral Higgs boson to the position of the zero value of the FBA. Our results show that the branching ratio of the process $B\to K\mu^+\mu^-$ can be quite largely modified by the effects of neutral Higgs bosons (NHBS) and the FBA would not vanish. The FBA for $B\to Kl^+l^- (l=\mu, \tau)$ vanishes if the contributions of NHB's vanish. The contributions of NHB's can be large enough to be observed only in SUSY and/or 2HDM with large $\tan\beta$, and a nonzero FBA for $B \to Kl^+l^-$ would signal the existance of new physics. For the process $B\to K^*\mu^+\mu^-$, the lepton transverse polarization is quite sensitive to the effects of NHBs, while the invariant mass spectrum, FBA, and lepton longitudinal polarization are not. For both $B\to K\tau^+\tau^-$ and $B\to K^*\tau^+\tau^-$, the effects of NHB's are quite significant. Our analysis manifests that even taking into account the theoretical uncertainties in calculating weak form factors, the effects of NHB's could bring SUSY to light. In brief, our analysis manifests that the effects of NHB's are quite remarkable in some regions of parameter space of SUSY, even for the process $B\to K\mu^+\mu^-$.

The paper is organized as follows. In Sec. II, the effective Hamiltonian is presented and the form factors given by using light cone sum rule method are briefly discussed. Basic formulas of observables are introduced in Sec. III. Section IV is devoted to the numerical analysis. In Sec. V we present discussions and conclusions.

II. EFFECTIVE HAMILTONIAN AND FORM FACTORS

By integrating out the degrees of heavy freedom from the full theory, MSSM, at the electroweak (EW) scale, we can get the effective Hamiltonian describing the rare semileptonic decay $b\to sl^+l^-$:

$$\mathcal{H}_{\rm eff}=-\frac{4G_F}{\sqrt{2}}V_{tb}V_{ts}^*\left(\sum_{i=1}^{10}C_i(\mu)O_i(\mu)+\sum_{i=1}^{10}C_{Q_i}(\mu)Q_i(\mu)\right), \tag{2.1}$$

where the first ten operators and Wilson coefficients at EW scale can be found in [8,38],[1] and the last ten operators and Wilson coefficients which represent the contributions of neutral Higgs boson can be found in [4].

With the renormalization-group equations to resum the QCD corrections, Wilson coefficients at energy scale $\mu=m_b$ are evaluated. Theoretical uncertainties related to renormalization scale can be substantially reduced when the next-leading-logarithm corrections are included [39].

The above Hamiltonian leads to the following free quark decay amplitude:

$$\begin{aligned}\mathcal{M}(b\to sl^+l^-)&=-\frac{G_F\alpha}{\sqrt{2}\pi}V_{ts}^*V_{tb}\Bigg\{C_9^{\rm eff}[\bar{s}\gamma_\mu Lb][\bar{l}\gamma^\mu l]\\&+C_{10}[\bar{s}\gamma_\mu Lb][\bar{l}\gamma^\mu\gamma_5 l]-2\hat{m}_bC_7^{\rm eff}\left[\bar{s}i\sigma_{\mu\nu}\frac{\hat{q}^\nu}{\hat{s}}Rb\right]\\&\times[\bar{l}\gamma^\mu l]+C_{Q1}[\bar{s}Rb][\bar{l}l]+C_{Q2}[\bar{s}Rb][\bar{l}\gamma_5 l]\Bigg\},\end{aligned} \tag{2.2}$$

where C_9^{eff} is defined as [40,41]

$$C_9^{\rm eff}(\mu,\hat{s})=C_9(\mu)+Y(\mu,\hat{s})+\frac{3\pi}{\alpha^2}C(\mu)\sum_{V_i=\psi(1s),\dots,\psi(6s)}\times\kappa_i\frac{\Gamma(V_i\to l^+l^-)\,m_{V_i}}{m_{V_i}{}^2-\hat{s}\,m_B^2-im_{V_i}\Gamma_{V_i}} \tag{2.3}$$

where $\hat{s}=s/m_b^2$, $s=q^2$, $C(\mu)=(3\,C_1+C_2+3\,C_3+C_4+3\,C_5+C_6)$, and

[1] In our previous papers, e.g., [3,4], we follow the convention of Ref. [1] for the indices of operators as well as Wilson coefficients. In this paper, we use more popular conventions (see, e.g., [39]). That is, $O_8\to O_9$ and $O_9\to O_{10}$.

$$Y(\mu,\hat{s})=g(\hat{m}_c,\hat{s})C(\mu)-\frac{1}{2}g(1,\hat{s})(4\,C_3+4\,C_4+3\,C_5+C_6)-\frac{1}{2}g(0,\hat{s})(C_3+3\,C_4)-\frac{2}{9}(3\,C_3+C_4+3\,C_5+C_6), \tag{2.4}$$

where the function $g(\hat{m}_c,\hat{s})$ comes from one-loop contributions of four-quark operators and is defined by

$$g(z,\hat{s})=-\frac{4}{9}\ln z^2+\frac{8}{27}+\frac{16}{9}\frac{z^2}{\hat{s}}-\begin{cases}\frac{2}{9}\sqrt{1-\frac{4z^2}{\hat{s}}}\left(2+\frac{4z^2}{\hat{s}}\right)\left[\ln\left(\frac{1+\sqrt{1-4z^2/\hat{s}}}{1-\sqrt{1-4z^2/\hat{s}}}\right)-i\pi\right], & 4z^2<\hat{s}\\ \frac{4}{9}\sqrt{\frac{4z^2}{\hat{s}}-1}\left(2+\frac{4z^2}{\hat{s}}\right)\arctan\left(\frac{1}{\sqrt{4z^2/\hat{s}-1}}\right), & 4z^2>\hat{s}.\end{cases} \tag{2.5}$$

The last terms in Eq. (2.3) are nonperturbative effects from $(\bar{c}c)$ resonance contributions, while the phenomenological factors κ_i can be fixed from the processes [23] $B\to K^{(*)}V_i\to K^{(*)}l^+l^-$ and as given in the Table I.

Exclusive decays $B\to(K,K^*)l^+l^-$ are described in terms of matrix elements of the quark operators in Eq. (2.2) over meson states, which can be parametrized in terms of form factors.

For the process $B\to Kl^+l^-$, the nonvanishing matrix elements are $(q=p_B-p)$

$$\langle K(p)|\bar{s}\gamma_\mu b|B(p_B)\rangle=f_+(s)\left\{(p_B+p)_\mu-\frac{m_B^2-m_K^2}{s}q_\mu\right\}+\frac{m_B^2-m_K^2}{s}f_0(s)\,q_\mu, \tag{2.6}$$

and

$$\begin{aligned}&\langle K(p)|\bar{s}\sigma_{\mu\nu}q^\nu(1+\gamma_5)b|B(p_B)\rangle\\&=\langle K(p)|\bar{s}\sigma_{\mu\nu}q^\nu b|B(p_B)\rangle\\&=i\{(p_B+p)_\mu s-q_\mu(m_B^2-m_K^2)\}\frac{f_T(s)}{m_B+m_K}.\end{aligned} \tag{2.7}$$

While for $B\to K^*l^+l^-$, related transition-matrix elements are

$$\begin{aligned}&\langle K^*(p)|(V-A)_\mu|B(p_B)\rangle\\&=-i\epsilon_\mu^*(m_B+m_{K^*})A_1(s)\\&\quad+i(p_B+p)_\mu(\epsilon^*p_B)\frac{A_2(s)}{m_B+m_{K^*}}\\&\quad+iq_\mu(\epsilon^*p_B)\frac{2m_{K^*}}{s}[A_3(s)-A_0(s)]\\&\quad+\epsilon_{\mu\nu\rho\sigma}\epsilon^{*\nu}p_B^\rho p^\sigma\frac{2V(s)}{m_B+m_{K^*}}.\end{aligned} \tag{2.8}$$

and

$$\begin{aligned}&\langle K^*|\bar{s}\sigma_{\mu\nu}q^\nu(1+\gamma_5)b|B(p_B)\rangle\\&=i\epsilon_{\mu\nu\rho\sigma}\epsilon^{*\nu}p_B^\rho p^\sigma\,2T_1(s)+T_2(s)\{\epsilon_\mu^*(m_B^2-m_{K^*}^2)\\&\quad-(\epsilon^*p_B)(p_B+p)_\mu\}+T_3(s)(\epsilon^*p_B)\\&\quad\times\left\{q_\mu-\frac{s}{m_B^2-m_{K^*}^2}(p_B+p)_\mu\right\},\end{aligned} \tag{2.9}$$

where ϵ_μ is the polarization vector of the vector meson K^*. By means of the equation of motion, one obtains several relations between form factors

$$\begin{aligned}&A_3(s)=\frac{m_B+m_{K^*}}{2m_{K^*}}A_1(s)-\frac{m_B-m_{K^*}}{2m_{K^*}}A_2(s),\\&A_0(0)=A_3(0),\\&\langle K^*|\partial_\mu A^\mu|B\rangle=2m_{K^*}(\epsilon^*p_B)A_0(s),\\&T_1(0)=T_2(0).\end{aligned} \tag{2.10}$$

All signs are defined in such a way as to render the form factors real and positive. The physical range in $\hat{s}$ extends from $\hat{s}_{\min}=4\hat{m}_l^2$ to $\hat{s}_{\max}=(1-\hat{m}_{K,K^*})^2$.

The calculation of the form factors given above is a real task, and one has to rely on certain approximate methods. We use the results calculated by using the technique of LCSR's and given in [23]. The form factors can be parametrized as

$$F(\hat{s})=F(0)\exp(c_1\hat{s}+c_2\hat{s}^2+c_3\hat{s}^3). \tag{2.11}$$

TABLE I. Fudge factors in $B\to K^{(*)}J/\Psi,\Psi'\to K^{(*)}l^+l^-$ decays calculated using the LCSR form factors.

κ	J/Ψ	Ψ'
K	2.70	3.51
K^*	1.65	2.36

The parametrization formula works within 1% accuracy for $s<15\ \mathrm{GeV}^2$ and can avoid the spurious singularities at $s=m_B^2$. Related parameters are given in Table IV of [23].

III. FORMULA OF OBSERVABLES

In this section we provide a formula for experimental observables, which include dilepton invariant mass spectrum, FBA, and lepton polarizations.

From Eqs. (2.2)–(2.8), it is straightforward to obtain the matrix element of $B\to K(K^*)l^+l^-$ as follows:

$$\mathcal{M}=-\frac{G_F\alpha}{2\sqrt{2}\pi}V_{ts}^*V_{tb}m_B \times[\mathcal{T}_\mu^1(\bar{l}\gamma^\mu l)+\mathcal{T}_\mu^2(\bar{l}\gamma^\mu\gamma_5 l)+\mathcal{S}(\bar{l}l)], \quad (3.1)$$

where for $B\to Kl^+l^-$,

$$\mathcal{T}_\mu^1=A'(\hat{s})\,\hat{p}_\mu, \quad (3.2)$$

$$\mathcal{T}_\mu^2=C'(\hat{s})\,\hat{p}_\mu+D'(\hat{s})\,\hat{q}_\mu, \quad (3.3)$$

$$\mathcal{S}=\mathcal{S}_1(\hat{s}), \quad (3.4)$$

and for $B\to K^*l^+l^-$,

$$\mathcal{T}_\mu^1=A(\hat{s})\,\epsilon_{\mu\rho\alpha\beta}\epsilon^{*\rho}\hat{p}_B^\alpha\hat{p}_{K^*}^\beta-iB(\hat{s})\,\epsilon^*{}_\mu+iC(\hat{s})\,(\epsilon^*\cdot\hat{p}_B)\hat{p}_\mu, \quad (3.5)$$

$$\mathcal{T}_\mu^2=E(\hat{s})\,\epsilon_{\mu\rho\alpha\beta}\epsilon^{*\rho}\hat{p}_B^\alpha\hat{p}_{K^*}^\beta-iF(\hat{s})\,\epsilon^*{}_\mu+iG(\hat{s})\,(\epsilon^*\cdot\hat{p}_B)\hat{p}_\mu +iH(\hat{s})\,(\epsilon^*\cdot\hat{p}_B)\hat{q}_\mu, \quad (3.6)$$

$$\mathcal{S}=i2\hat{m}_{K^*}(\epsilon^*\cdot\hat{p}_B)\mathcal{S}_2(\hat{s}) \quad (3.7)$$

with $p\equiv p_B+p_{K,K^*}$. Note that, using the equation of motion for lepton fields, the terms in $\hat{q}_\mu$ in $\mathcal{T}_\mu^1$ vanish.

The auxiliary functions above are defined as

$$A'(\hat{s})=C_9{}^{\mathrm{eff}}(\hat{s})\,f_+(\hat{s})+\frac{2\hat{m}_b}{1+\hat{m}_K}C_7{}^{\mathrm{eff}}f_T(\hat{s}), \quad (3.8)$$

$$C'(\hat{s})=C_{10}\,f_+(\hat{s}), \quad (3.9)$$

$$D'(\hat{s})=C_{10}\,f_-(\hat{s})-\frac{1-\hat{m}_K^2}{2\hat{m}_l(\hat{m}_b-\hat{m}_s)}C_{Q2}f_0(\hat{s}), \quad (3.10)$$

$$\mathcal{S}_1(\hat{s})=\frac{1-\hat{m}_K^2}{(\hat{m}_b-\hat{m}_s)}C_{Q1}f_0(\hat{s}), \quad (3.11)$$

$$A(\hat{s})=\frac{2}{1+\hat{m}_{K^*}}C_9{}^{\mathrm{eff}}(\hat{s})V(\hat{s})+\frac{4\hat{m}_b}{\hat{s}}C_7{}^{\mathrm{eff}}T_1(\hat{s}), \quad (3.12)$$

$$B(\hat{s})=(1+\hat{m}_{K^*})\left[C_9{}^{\mathrm{eff}}(\hat{s})A_1(\hat{s})+\frac{2\hat{m}_b}{\hat{s}}\times(1-\hat{m}_{K^*})C_7{}^{\mathrm{eff}}T_2(\hat{s})\right], \quad (3.13)$$

$$C(\hat{s})=\frac{1}{1-\hat{m}_{K^*}^2}\left[(1-\hat{m}_{K^*})C_9{}^{\mathrm{eff}}(\hat{s})A_2(\hat{s}) +2\hat{m}_bC_7{}^{\mathrm{eff}}\left(T_3(\hat{s})+\frac{1-\hat{m}_{K^*}^2}{\hat{s}}T_2(\hat{s})\right)\right], \quad (3.14)$$

$$E(\hat{s})=\frac{2}{1+\hat{m}_{K^*}}C_{10}V(\hat{s}), \quad (3.15)$$

$$F(\hat{s})=(1+\hat{m}_{K^*})C_{10}A_1(\hat{s}), \quad (3.16)$$

$$G(\hat{s})=\frac{1}{1+\hat{m}_{K^*}}C_{10}A_2(\hat{s}), \quad (3.17)$$

$$H(\hat{s})=\frac{C_{10}}{\hat{s}}[(1+\hat{m}_{K^*})A_1(\hat{s})-(1-\hat{m}_{K^*})A_2(\hat{s}) -2\hat{m}_{K^*}A_0(\hat{s})]+\frac{\hat{m}_{K^*}}{\hat{m}_l(\hat{m}_b+\hat{m}_s)}A_0(\hat{s})C_{Q2}, \quad (3.18)$$

$$\mathcal{S}_2(\hat{s})=-\frac{1}{(\hat{m}_b+\hat{m}_s)}A_0(\hat{s})C_{Q1}, \quad (3.19)$$

where

$$f_0(\hat{s})=\frac{1}{1-\hat{m}_K^2}[\hat{s}f_-(\hat{s})+(1-\hat{m}_K^2)f_+(\hat{s})]. \quad (3.20)$$

To get the auxiliary functions given above, we have used the equations of motion

$$q^\mu(\bar{\psi}_1\gamma_\mu\psi_2)=(m_1-m_2)\bar{\psi}_1\psi_2, \quad (3.21)$$

$$q^\mu(\bar{\psi}_1\gamma_\mu\gamma_5\psi_2)=-(m_1+m_2)\bar{\psi}_1\gamma_5\psi_2. \quad (3.22)$$

The contributions of NHB's have been incorporated in the terms of $\mathcal{S}_1(\hat{s})$, $D'(\hat{s})$, $H(\hat{s})$, and $\mathcal{S}_2(\hat{s})$. It is remarkable that the contributions of NHB's in $D'(\hat{s})$ and $H(\hat{s})$ are proportional to the inverse mass of the lepton, and for the case

$l=\mu$, the effects of NHB's can be manifested through these terms.

A phenomenological effective Hamiltonian was recently given in [30]. If we ignore tensor-type interactions in the phenomenological Hamiltonian (it is shown that physical observables are not sensitive to the presence of tensor-type interactions [6]), it is easy to verify that the matrix element of $B\to K^{(*)}l^+l^-$ can always be expressed as the form of Eq. (3.1) with the auxiliary functions defined as

$$A'(\hat{s})=w_{c_1}f_+(\hat{s})-\frac{w_{c_9}+w_{c_{10}}}{1+\hat{m}_K}f_T(\hat{s}), \tag{3.23}$$

$$C'(\hat{s})=w_{c_2}f_+(\hat{s}), \tag{3.24}$$

$$D'(\hat{s})=w_{c_2}f_-(\hat{s})-\frac{1-\hat{m}_K^2}{2\hat{m}_l(\hat{m}_b-\hat{m}_s)}w_{c_6}f_0(\hat{s}), \tag{3.25}$$

$$\mathcal{S}_1(\hat{s})=\frac{1-\hat{m}_K^2}{(\hat{m}_b-\hat{m}_s)}w_{c_5}f_0(\hat{s}), \tag{3.26}$$

$$A(\hat{s})=\frac{2}{1+\hat{m}_{K^*}}w_{c_1}V(\hat{s})-\frac{2}{\hat{s}}(w_{c_9}+w_{c_{10}})T_1(\hat{s}), \tag{3.27}$$

$$B(\hat{s})=-(1+\hat{m}_{K^*})\left[w_{c_3}A_1(\hat{s})+\frac{1}{\hat{s}}(1-\hat{m}_{K^*})(w_{c_9}+w_{c_{10}})T_2(\hat{s})\right], \tag{3.28}$$

$$C(\hat{s})=-\frac{1}{1-\hat{m}_{K^*}^2}\left[(1-\hat{m}_{K^*})w_{c_3}(\hat{s})A_2(\hat{s})+(w_{c_9}-w_{c_{10}})\left((1+\hat{m}_{K^*})T_3(\hat{s})+\frac{1-\hat{m}_{K^*}^2}{\hat{s}}T_2(\hat{s})\right)\right], \tag{3.29}$$

$$E(\hat{s})=\frac{2}{1+\hat{m}_{K^*}}w_{c_2}V(\hat{s}), \tag{3.30}$$

$$F(\hat{s})=-(1+\hat{m}_{K^*})w_{c_4}A_1(\hat{s}), \tag{3.31}$$

$$G(\hat{s})=-\frac{1}{1+\hat{m}_{K^*}}w_{c_4}A_2(\hat{s}), \tag{3.32}$$

$$H(\hat{s})=-\frac{2\hat{m}_{K^*}}{\hat{s}}w_{c_4}[A_3(\hat{s})-A_0(\hat{s})]+\frac{\hat{m}_{K^*}}{\hat{m}_l(\hat{m}_b+\hat{m}_s)}w_{c_8}A_0(\hat{s}), \tag{3.33}$$

$$\mathcal{S}_2(\hat{s})=-\frac{1}{(\hat{m}_b+\hat{m}_s)}w_{c_7}A_0(\hat{s}), \tag{3.34}$$

where

$$w_{c_1}=\frac{1}{4}(C_{LL}+C_{LR}+C_{RL}+C_{RR}), \tag{3.35}$$

$$w_{c_2}=\frac{1}{4}(-C_{LL}+C_{LR}-C_{RL}+C_{RR}), \tag{3.36}$$

$$w_{c_3}=\frac{1}{4}(-C_{LL}-C_{LR}+C_{RL}+C_{RR}), \tag{3.37}$$

$$w_{c_4}=\frac{1}{4}(C_{LL}-C_{LR}-C_{RL}+C_{RR}), \tag{3.38}$$

$$w_{c_5}=\frac{1}{4}(C_{LRLR}+C_{RLLR}+C_{LRRL}+C_{RLRL}), \tag{3.39}$$

$$w_{c_6}=\frac{1}{4}(C_{LRLR}+C_{RLLR}-C_{LRRL}-C_{RLRL}), \tag{3.40}$$

$$w_{c_7}=\frac{1}{4}(C_{LRLR}-C_{RLLR}+C_{LRRL}-C_{RLRL}), \tag{3.41}$$

$$w_{c_8}=\frac{1}{4}(C_{LRLR}-C_{RLLR}-C_{LRRL}+C_{RLRL}), \tag{3.42}$$

$$w_{c_9}=m_bC_{BR}, \tag{3.43}$$

$$w_{c_{10}}=m_sC_{SL}. \tag{3.44}$$

In the above equations C_{LL},C_{LR}, etc., are defined in Ref. [6]. Therefore, our formula given below can also be used to make model-independent phenomenological analysis, if using Eqs. (3.23)–(3.34) instead of Eqs. (3.8)–(3.19).

Keeping the lepton mass, we find the double differential decay widths Γ^K and Γ^{K^*} for the decays $B\to Kl^+l^-$ and $B\to K^*l^+l^-$, respectively, as

$$\frac{d^2\Gamma^K}{d\hat{s}d\hat{u}}=\frac{G_F^2\alpha^2 m_B^5}{2^{11}\pi^5}|V_{ts}^* V_{tb}|^2\{(|A'|^2+|C'|^2)(\lambda-\hat{u}^2)+|\mathcal{S}_1|^2(\hat{s}-4\hat{m}_l^2)+\mathrm{Re}(\mathcal{S}_1 A'^*)4\hat{m}_l\hat{u}+|C'|^2 4\hat{m}_l^2(2+2\hat{m}_K^2-\hat{s})$$
$$+\mathrm{Re}(C'D'^*)8\hat{m}_l^2(1-\hat{m}_K^2)+|D'|^2 4\hat{m}_l^2\hat{s}\}, \tag{3.45}$$

$$\frac{d^2\Gamma^{K^*}}{d\hat{s}d\hat{u}}=\frac{G_F^2\,\alpha^2\,m_B^5}{2^{11}\pi^5}|V_{ts}^*\,V_{tb}|^2\left\{\frac{|A|^2}{4}[\hat{s}(\lambda+\hat{u}^2)+4\hat{m}_l^2\lambda]+\frac{|E|^2}{4}[\hat{s}(\lambda+\hat{u}^2)-4\hat{m}_l^2\lambda]+|\mathcal{S}_2|^2(\hat{s}-4\hat{m}_l^2)\lambda\right.$$
$$+\frac{1}{4\hat{m}_{K*}^2}[|B|^2(\lambda-\hat{u}^2+8\hat{m}_{K*}^2(\hat{s}+2\hat{m}_l^2))+|F|^2(\lambda-\hat{u}^2+8\hat{m}_{K*}^2(\hat{s}-4\hat{m}_l^2))]-2\hat{s}\hat{u}[\mathrm{Re}(BE^*)+\mathrm{Re}(AF^*)]$$
$$+\frac{2\hat{m}_l\hat{u}}{\hat{m}_{K*}}[\mathrm{Re}(\mathcal{S}_2B^*)(\hat{s}+\hat{m}_{K*}^2-1)+\mathrm{Re}(\mathcal{S}_2C^*)\lambda]+\frac{\lambda}{4\hat{m}_{K*}^2}[|C|^2(\lambda-\hat{u}^2)+|G|^2(\lambda-\hat{u}^2+4\hat{m}_l^2(2+2\hat{m}_{K*}^2-\hat{s}))]$$
$$-\frac{1}{2\hat{m}_{K*}^2}[\mathrm{Re}(BC^*)(1-\hat{m}_{K*}^2-\hat{s})(\lambda-\hat{u}^2)+\mathrm{Re}(FG^*)((1-\hat{m}_{K*}^2-\hat{s})(\lambda-\hat{u}^2)+4\hat{m}_l^2\lambda)]-2\frac{\hat{m}_l^2}{\hat{m}_{K*}^2}\lambda[\mathrm{Re}(FH^*)$$
$$\left.-\mathrm{Re}(GH^*)(1-\hat{m}_{K*}^2)]+|H|^2\frac{\hat{m}_l^2}{\hat{m}_{K*}^2}\hat{s}\lambda\right\}. \tag{3.46}$$

Here the kinematic variables $(\hat{s},\hat{u})$ are defined as

$$\hat{s}=\hat{q}^2=(\hat{p}_+ +\hat{p}_-)^2, \tag{3.47}$$

$$\hat{u}=(\hat{p}_B-\hat{p}_-)^2-(\hat{p}_B-\hat{p}_+)^2 \tag{3.48}$$

which are bounded as

$$(2\hat{m}_l)^2\leqslant\hat{s}\leqslant(1-\hat{m}_{K,K*})^2, \tag{3.49}$$

$$-\hat{u}(\hat{s})\leqslant\hat{u}\leqslant\hat{u}(\hat{s}), \tag{3.50}$$

with $\hat{m}_l=m_l/m_B$ and

$$\hat{u}(\hat{s})=\sqrt{\lambda\left(1-4\frac{\hat{m}_l^2}{\hat{s}}\right)}, \tag{3.51}$$

$$\lambda=1+\hat{m}_{K,K*}^4+\hat{s}^2-2\hat{s}-2\hat{m}_{K,K*}^2(1+\hat{s}), \tag{3.52}$$

$$\mathcal{D}=\sqrt{1-\frac{4\hat{m}_l^2}{s}}. \tag{3.53}$$

Note that the variable $\hat{u}$ corresponds to θ, the angle between the momentum of the B meson and the positively charged lepton l^+ in the dilepton c.m. system (c.m.s.) frame, through the relation $\hat{u}=-\hat{u}(\hat{s})\cos\theta$ [42].

Integrating over $\hat{u}$ in the kinematic region given in Eq. (3.50) we get the formula of dilepton invariant mass spectra (IMS)

$$\frac{d\Gamma^K}{d\hat{s}}=\frac{G_F^2\alpha^2 m_B^5}{2^{10}\pi^5}|V_{ts}^*V_{tb}|^2\hat{u}(\hat{s})D^K \tag{3.54}$$

$$D^K=(|A'|^2+|C'|^2)\left(\lambda-\frac{\hat{u}(\hat{s})^2}{3}\right)+|\mathcal{S}_1|^2(\hat{s}-4\hat{m}_l^2)$$
$$+|C'|^2 4\hat{m}_l^2(2+2\hat{m}_K^2-\hat{s})+\mathrm{Re}(C'D''^*)8\hat{m}_l^2(1-\hat{m}_K^2)+|D'|^2 4\hat{m}_l^2\hat{s}, \tag{3.55}$$

$$\frac{d\Gamma^{K^*}}{d\hat{s}}=\frac{G_F^2\,\alpha^2\,m_B^5}{2^{10}\pi^5}|V_{ts}^{*}V_{tb}|^2\,\hat{u}(\hat{s})D^{K^*} \tag{3.56}$$

$$\begin{aligned}D^{K^*}=&\frac{|A|^2}{3}\hat{s}\lambda\left(1+2\frac{\hat{m}_l^2}{\hat{s}}\right)+|E|^2\hat{s}\frac{\hat{u}(\hat{s})^2}{3}+|\mathcal{S}_2|^2(\hat{s}-4\hat{m}_l^2)\lambda+\frac{1}{4\hat{m}_{K^*}^2}\left[|B|^2\left(\lambda-\frac{\hat{u}(\hat{s})^2}{3}+8\hat{m}_{K^*}^2(\hat{s}+2\hat{m}_l^2)\right)\right.\\&\left.+|F|^2\left(\lambda-\frac{\hat{u}(\hat{s})^2}{3}+8\hat{m}_{K^*}^2(\hat{s}-4\hat{m}_l^2)\right)\right]+\frac{\lambda}{4\hat{m}_{K^*}^2}\left[|C|^2\left(\lambda-\frac{\hat{u}(\hat{s})^2}{3}\right)+|G|^2\left(\lambda-\frac{\hat{u}(\hat{s})^2}{3}\right.\right.\\&\left.\left.+4\hat{m}_l^2(2+2\hat{m}_{K^*}^2-\hat{s})\right)\right]-\frac{1}{2\hat{m}_{K^*}^2}\left[\mathrm{Re}(BC^*)\left(\lambda-\frac{\hat{u}(\hat{s})^2}{3}\right)(1-\hat{m}_{K^*}^2-\hat{s})\right.\\&\left.+\mathrm{Re}(FG^*)\left(\left(\lambda-\frac{\hat{u}(\hat{s})^2}{3}\right)(1-\hat{m}_{K^*}^2-\hat{s})+4\hat{m}_l^2\lambda\right)\right]-2\frac{\hat{m}_l^2}{\hat{m}_{K^*}^2}\lambda[\mathrm{Re}(FH^*)-\mathrm{Re}(GH^*)(1-\hat{m}_{K^*}^2)]\\&+\frac{\hat{m}_l^2}{\hat{m}_{K^*}^2}\hat{s}\lambda|H|^2.\end{aligned} \tag{3.57}$$

Both distributions agree with the ones obtained in [23,36], if $C_{Q_1,2}$ are set to zero.

The differential FBA is defined as

$$A_{FB}(s)=\frac{-\int_0^{u(\hat{s})}dz\,\frac{d\Gamma}{dsdu}+\int_{-u(\hat{s})}^{0}du\frac{d\Gamma}{dsdu}}{\int_0^{u(\hat{s})}dz\,\frac{d\Gamma}{dsdu}+\int_{-u(\hat{s})}^{0}du\,\frac{d\Gamma}{dsdu}}.$$

For $B\to Kl^+l^-$ decays it reads as follows:

$$\frac{d\mathcal{A}_{\mathrm{FB}}^{K}}{d\hat{s}}D^K=-2\hat{m}_l\hat{u}(\hat{s})\mathrm{Re}(\mathcal{S}_1A'^*). \tag{3.58}$$

For $B\to K^*l^+l^-$ decays it reads as follows:

$$\begin{aligned}\frac{d\mathcal{A}_{\mathrm{FB}}^{K^*}}{d\hat{s}}D^{K^*}=\hat{u}(\hat{s})\Big\{&\hat{s}[\mathrm{Re}(BE^*)+\mathrm{Re}(AF^*)]\\&+\frac{\hat{m}_l}{\hat{m}_{K^*}}[\mathrm{Re}(\mathcal{S}_2B^*)(1-\hat{s}-\hat{m}_{K^*}^2)\\&-\mathrm{Re}(\mathcal{S}_2C^*)\lambda]\Big\}.\end{aligned} \tag{3.59}$$

We can read from Eq. (3.58), the FBA of the process $B\to Kl^+l^-$ does not vanish when the contributions of NHB are taken into account. With it, our analysis below also shows that the contributions of NHB's can even be accessible in B factories.

The lepton polarization can be defined as follows:

$$\frac{d\Gamma(\vec{n})}{ds}=\frac{1}{2}\left(\frac{d\Gamma}{ds}\right)_0[1+(P_L\vec{e}_L+P_N\vec{e}_N+P_T\vec{e}_T)\cdot\vec{n}], \tag{3.60}$$

where the subscript ''0'' corresponds to the unpolarized width, and P_L, P_T, and P_N, correspond to the longitudinal, transverse, and normal components of the polarization vector, respectively.

TABLE II. Values of the input parameters used in our numerical analysis.

m_b	4.8 GeV
m_c	1.4 GeV
m_s	0.2 GeV
m_μ	0.11 GeV
m_τ	1.78 GeV
M_B	5.28 GeV
M_K	0.49 GeV
M_{K^*}	0.89 GeV
$M_{J/\psi}(M_{\psi'})$	3.10(3.69) GeV
Γ_B	4.22×10^{-13} GeV
$\Gamma_{J/\psi}(\Gamma_{\psi'})$	$8.70(27.70)\times10^{-5}$ GeV
$\Gamma(J/\psi\to l^+l^-)$	5.26×10^{-6} GeV
$\Gamma(\psi'\to l^+l^-)$	2.14×10^{-6} GeV
G_F	1.17×10^{-5} GeV^{-2}
α^{-1}	129
$\lvert V_{ts}^*V_{tb}\rvert$	0.0385

TABLE III. Wilson coefficients of the SM used in the numerical analysis.

C_1	C_2	C_3	C_4	C_5	C_6	$C_7^{\rm eff}$	C_9	C_{10}	C
-0.248	$+1.107$	$+0.011$	-0.026	$+0.007$	-0.031	-0.313	$+4.344$	-4.669	$+0.362$

For the process $B\to Kl^-l^+$, the P_L^K, P_T^K, and P_N^K, are derived, respectively, as

$$P_L^K D^K=\frac{4}{3}\mathcal{D}\{\lambda\,\mathrm{Re}(A'C'^*)-3\hat{m}_l(1-\hat{m}_K^2)\mathrm{Re}(C'^*\mathcal{S}_1)-3\hat{m}_l\hat{s}\mathrm{Re}(D'^*\mathcal{S}_1)\},\tag{3.61}$$

$$P_N^K D^K=\frac{\pi\sqrt{\hat{s}}\hat{u}(\hat{s})}{2}\{-\mathrm{Im}(A'\mathcal{S}_1^*)+2\hat{m}_l\,\mathrm{Im}(C'D'^*)\},\tag{3.62}$$

$$P_T^K D^K=\frac{-\pi\sqrt{\lambda}}{\sqrt{\hat{s}}}\left\{\hat{m}_l[(1-\hat{m}_K^2)\mathrm{Re}(A'C'^*)+\hat{s}\mathrm{Re}(A'D'^*)]+\frac{(\hat{s}-4\hat{m}_l^2)}{2}\mathrm{Re}(C'\mathcal{S}_1^*)\right\}.\tag{3.63}$$

D^K is defined in Eq. (3.55). For the process $B\to K^*l^-l^+$, the $P_L^{K^*}$, $P_T^{K^*}$, and $P_N^{K^*}$, are derived, respectively, as

$$P_L^{K^*}D^{K^*}=\mathcal{D}\left\{\frac{2\hat{s}\lambda}{3}\mathrm{Re}(AE^*)+\frac{(\lambda+12\hat{m}_{K^*}^2)}{3\hat{m}_{K^*}^2}\mathrm{Re}(BF^*)-\frac{\lambda(1-\hat{m}_{K^*}^2-\hat{s})}{3\hat{m}_{K^*}^2}\mathrm{Re}(BG^*+CF^*)+\frac{\lambda^2}{3\hat{m}_{K^*}^2}\mathrm{Re}(CG^*)+\frac{2\hat{m}_l\lambda}{\hat{m}_{K^*}}[\mathrm{Re}(F\mathcal{S}_2^*)-\hat{s}\mathrm{Re}(H\mathcal{S}_2^*)-(1-\hat{m}_{K^*}^2)\mathrm{Re}(G\mathcal{S}_2^*)]\right\},\tag{3.64}$$

$$P_N^{K^*}D^{K^*}=\frac{-\pi\sqrt{\hat{s}}\hat{u}(\hat{s})}{4\hat{m}_K}\left\{\frac{\hat{m}_l}{\hat{m}_{K^*}}[\mathrm{Im}(FG^*)(1+3\hat{m}_{K^*}^2-s)+\mathrm{Im}(FH^*)(1-\hat{m}_{K^*}^2-s)-\mathrm{Im}(GH^*)\lambda]+2\hat{m}_{K^*}\hat{m}_l[\mathrm{Im}(BE^*)+\mathrm{Im}(AF^*)]-(1-\hat{m}_{K^*}^2-\hat{s})\mathrm{Im}(B\mathcal{S}_2^*)+\lambda\,\mathrm{Im}(C\mathcal{S}_2^*)\right\},\tag{3.65}$$

$$P_T^{K^*}D^{K^*}=\frac{\pi\sqrt{\lambda}\hat{m}_l}{4\sqrt{\hat{s}}}\left\{4\hat{s}\mathrm{Re}(AB^*)+\frac{(1-\hat{m}_{K^*}^2-\hat{s})}{\hat{m}_{K^*}^2}[-\mathrm{Re}(BF^*)+(1-\hat{m}_{K^*}^2)\mathrm{Re}(BG^*)+\hat{s}\mathrm{Re}(BH^*)]+\frac{\lambda}{\hat{m}_{K^*}^2}[\mathrm{Re}(CF^*)-(1-\hat{m}_{K^*}^2)\mathrm{Re}(CG^*)-\hat{s}\mathrm{Re}(CH^*)]+\frac{(\hat{s}-4\hat{m}_l^2)}{\hat{m}_{K^*}\hat{m}_l}\times[(1-\hat{m}_{K^*}^2-\hat{s})\mathrm{Re}(F\mathcal{S}_2^*)-\lambda\,\mathrm{Re}(G\mathcal{S}_2^*)]\right\}.\tag{3.66}$$

D^{K^*} is defined by Eq. (3.57).

IV. NUMERICAL ANALYSIS

Parameters used in our analysis are listed in Table II. Considering that the branching ratios of $B\to Kl^+l^-$ and $B\to K^*l^+l^-$ are not very sensitive to the mass of m_b, we neglect the difference between the pole mass and running mass of the b quark.

The Wilson coefficients in the SM used in the numerical analysis are given in Table III. $C_7^{\rm eff}$ is defined as

$$C_7^{eff}=C_7-C_5/3-C_6.\tag{4.1}$$

TABLE IV. Wilson coefficients of the SUSY used in our numerical analysis. R_i means C_i/C_i^{SM}. SUSY I corresponds to the regions where SUSY can destructively contribute and can change the sign of C_7, but the contributions of NHB's are neglected. SUSY II corresponds to the regions where $\tan\beta$ is large and the masses of superpartners are relatively small. SUSY III corresponds to the regions where $\tan\beta$ is large but the masses of superpartners are relatively large. In the last two cases the effects of NHB's are taken into account. The contributions of NHB's are settled to be different for both the case $l=\mu$ and $l=\tau$, since $C_{Q_{1,2}}$ are proportional to the mass of lepton. The values in bracket are for the case $l=\tau$.

SUSY models	R_7	R_9	R_{10}	C_{Q_1}	C_{Q_2}
SUSY I	-1.2	1.1	0.8	0.0	0.0
SUSY II	-1.2	1.1	0.8	6.5(16.5)	$-6.5(-16.5)$
SUSY III	1.2	1.1	0.8	1.2(4.5)	$-1.2(-4.5)$

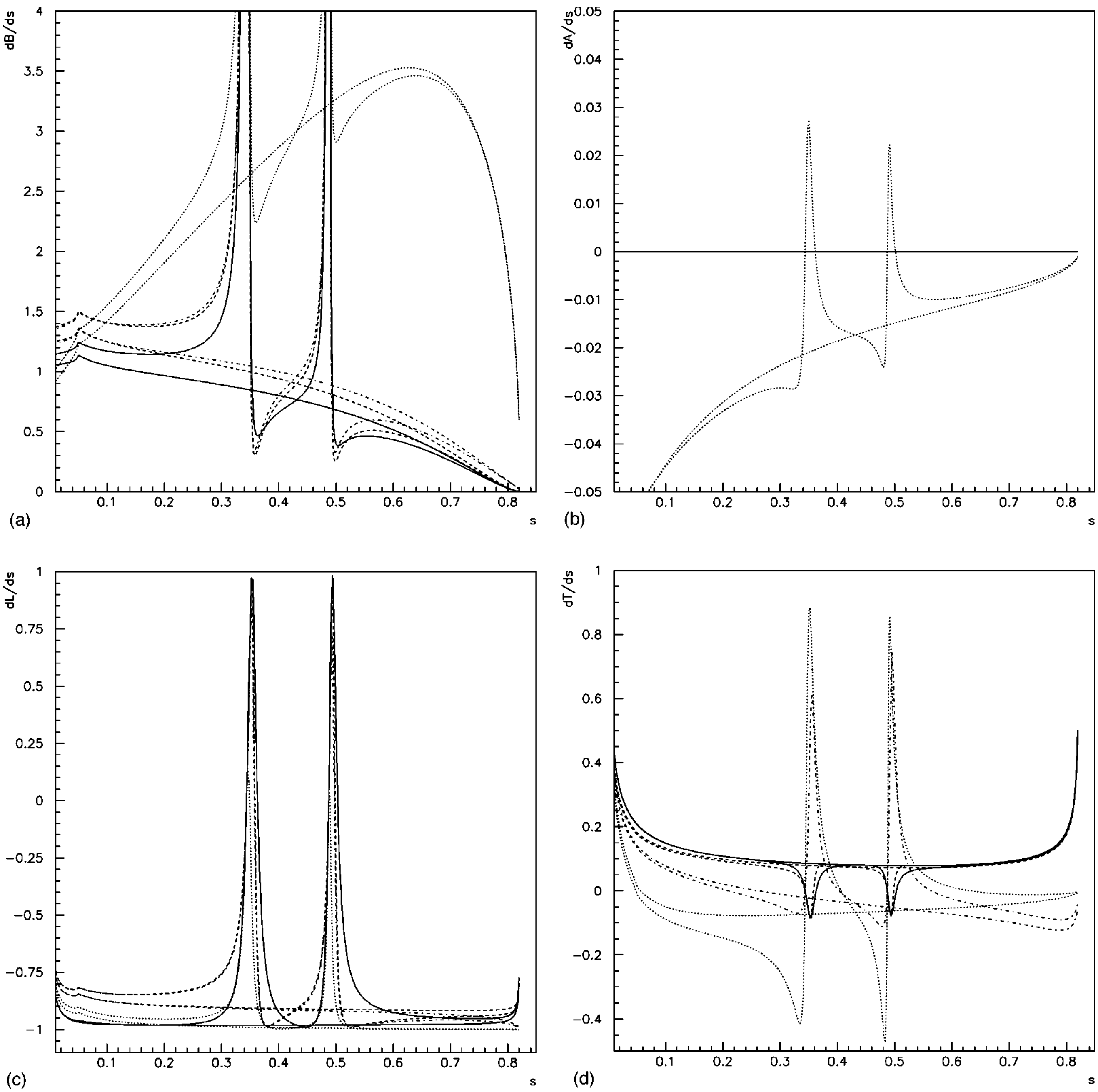

FIG. 1. The IMS(a), FBA(b), P_L(c), and P_T(d) of the process $B\to K\mu^+\mu^-$. The solid line, dashed line, dot line, and dashed-dot line represent the SM, SUSY I, SUSY II, SUSY III respectively. Both the total (SD+LD) and the pure SD contributions are shown in order to compare.

$C_{Q_{1,2}}$ come from exchanging NHB's and are proportional to $\tan^3\beta$ in some regions of the parameter space in SUSY models. According to the analysis in [4,37], the necessary conditions for the large contributions of NHB's include: (i) the ratio of vacuum expectation value, $\tan\beta$, should be large, (ii) the mass values of the lighter chargino and the lighter stop should not be too large (say less than 120 GeV), (iii) mass splitting of charginos and stops should be large, which also indicate large mixing between stop sector and chargino sector. As the conditions are satisfied, the process $B\to X_s\gamma$ will impose a constraint on C_7. It is well known that this process puts a very stringent constraint on the possible new physics and that SUSY can contribute destructively when the signature of the Higgs mass term μ is minus. There exist considerable regions of SUSY parameter space in which NHB's can largely contribute to the process $b\to sl^+l^-$, while the constraint of $b\to s\gamma$ is respected (i.e., the signature of the Wilson coefficient C_7 is changed from positive to negative). When the masses of the SUSY particle are relatively heavy (say, 450 GeV), there are still significant regions in the parameter space of SUSY models in which NHB's could contribute largely. However, at these cases C_7 does not change its sign, because contributions of charged Higgs and charginos cancel with each other. We will see it is

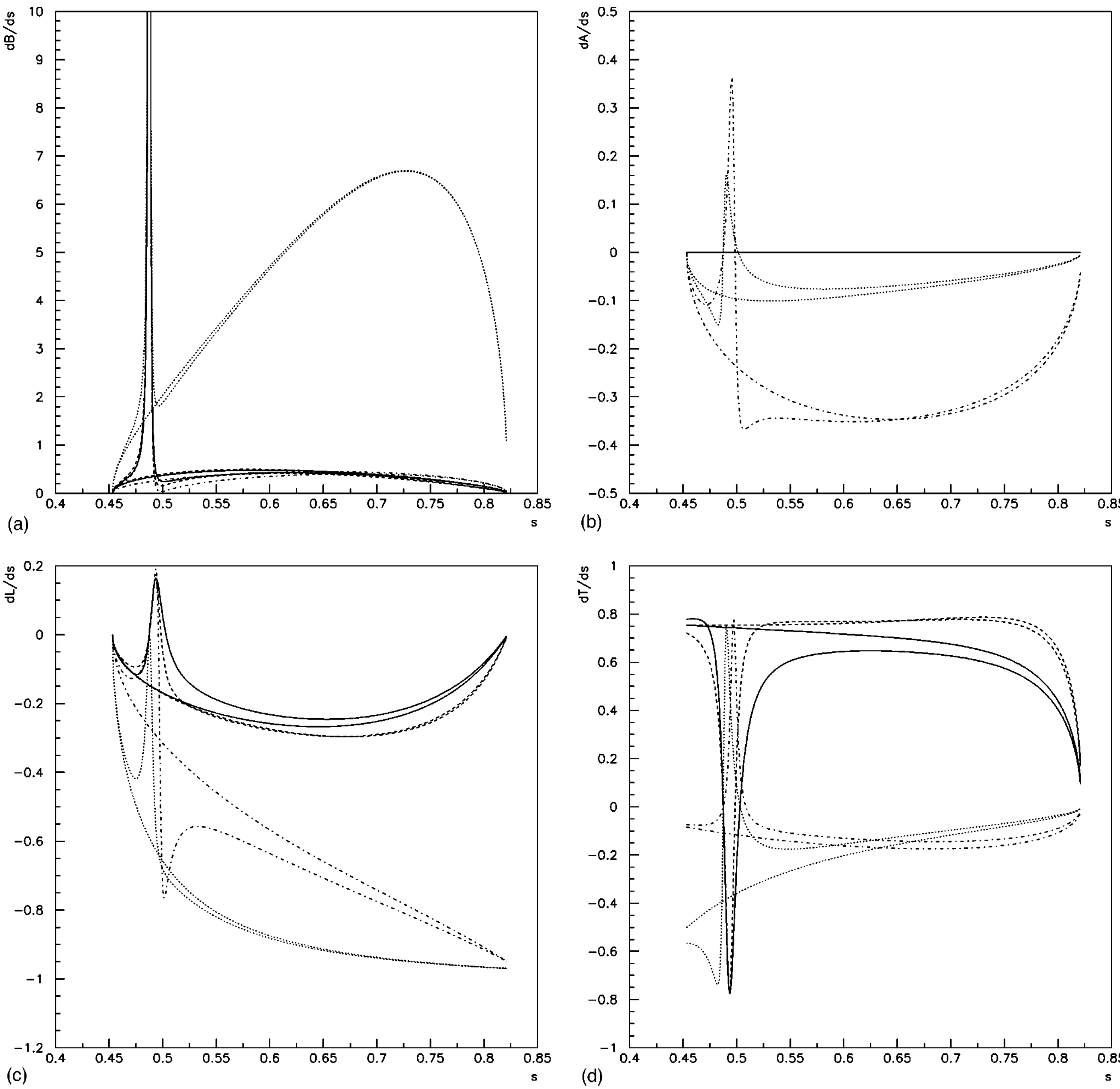

FIG. 2. The IMS(a), FBA(b), P_L(c), and P_T(d) of the process $B\rightarrow K\tau^+\tau^-$. The line conventions are the same as given in the legend of Fig. 1.

hopeful to distinguish these two kinds of regions of SUSY parameter space by observing $B\rightarrow K^{(*)}l^+l^-$.

As pointed out in [3,4], the contribution of NHB's is proportional to the lepton mass, therefore, for $l=e$, contributions of NHB's can be safely neglected, while for cases $l=\mu$ and $l=\tau$, the contributions of NHB's can be considerably large. To investigate the effects of NHB's in SUSY models, we take typical values of $C_{7,9,10}$ and $C_{Q_{1,2}}$ as given in Table IV. The SUSY model without considering the effects of NHB's (SUSY I in Table IV) is given as a reference frame so that the effects of NHB's could be shown in high relief.

Numerical results are shown in Figs. 1–4. In Fig. 1(a), the IMS of $B\rightarrow K\mu^+\mu^-$ is depicted. We see that at the high $\hat{s}$ regions, NHB's greatly modify the spectrum, while at the low $\hat{s}$ region, the effects of NHB's become weak. In Fig. 1(b), the FBA of the $B\rightarrow K\mu^+\mu^-$ is presented. Figure 1(b) shows that the average FBA in $B\rightarrow K\mu^+\mu^-$ is 0.02. To measure an asymmetry A of a decay with the branching ratio Br at the $n\sigma$ level, the required number of events is $N=n^2/(BrA^2)$. For $B\rightarrow K\mu^+\mu^-$, the average FBA is 0.02 or so, the required number of events is 10^{12} or so. Therefore, it is hard to observe the derivation of FBA from the SM. In Figs. 1(c) and 1(d), the longitudinal and transverse polarizations are given. The effect of NHB's on the longitudinal polarization is weak but the effect on the transverse is remarkable.

In Figs. 2(a) and 2(b) the IMS and FBA of $B\rightarrow K\tau^+\tau^-$

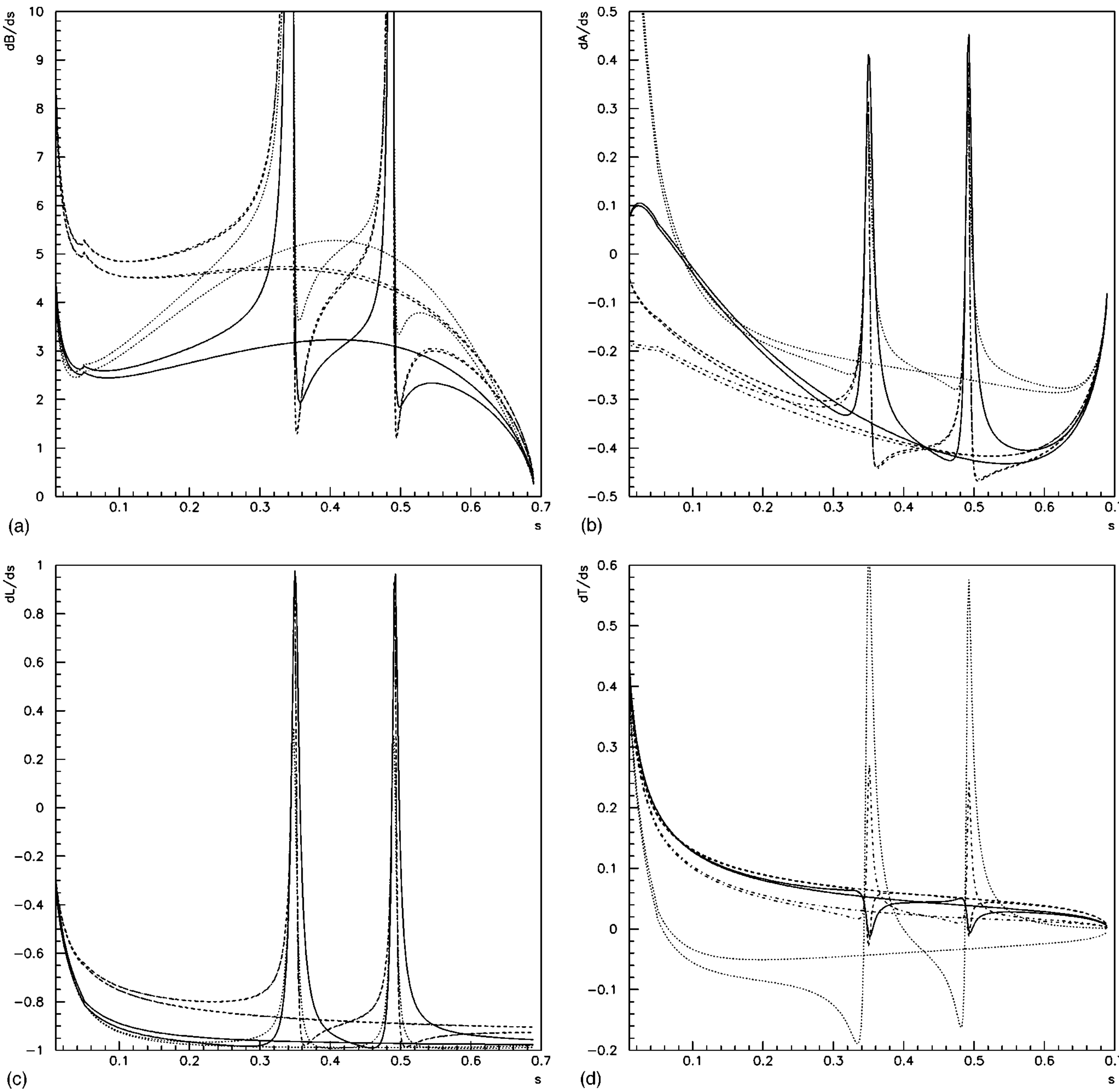

FIG. 3. The IMS(a), FBA(b), P_L(c), and P_T(d) of the process $B \to K^* \mu^+ \mu^-$. The line conventions are the same as given in the legend of Fig. 1.

are presented, respectively. For SUSY II, the effects of NHB's to IMS are quite manifest, and the average FBA can reach 0.1. For SUSY III, the average FBA can reach 0.3. Therefore, in order to observe FBA, the required number of events should be 10^9 or so and 10^8, respectively, so that in B factories, say LHCB, these two cases are accessible. In Figs. 2(c) and 2(d), the longitudinal and transverse polarizations are drawn, respectively. The effects of NHB's are also very obvious.

Figures 3 and 4 are devoted to the decay $B \to K^* l^+ l^-$. In Fig. 3, the IMS, FBA, and polarizations of $B \to K^* \mu^+ \mu^-$ are given. We see that this process is not as sensitive to the effect of NHB as $B \to K \mu^+ \mu^-$. However, the contribution of NHB's will increase the part with positive FBA and will be helpful to determine the zero point of FBA. Figure 3(d) depicts the transverse polarization of the $B \to K^* \mu^+ \mu^-$, and the effect of NHB's is quite obvious. The zero point of the FBA can be slightly modified as shown in Fig 3(b) due to the contributions of NHB's.

In Fig. 4, the IMS, FBA, longitudinal, and transverse polarizations of the $B \to K^* \tau^+ \tau^-$ are depicted. The effect of NHB's does show in great relief. It is worth noting that IMS, FBA, and lepton polarizations for $B \to K^* l^+ l^-$ in MSSM without including the contributions of NHB's are also significantly different from those in SM, while for $B \to K l^+ l^-$ they have little difference from those in SM. Therefore, compared to the process $B \to K l^+ l^-$, more precise measurements for $B \to K^* l^+ l^-$ are needed in order to single out the contri-

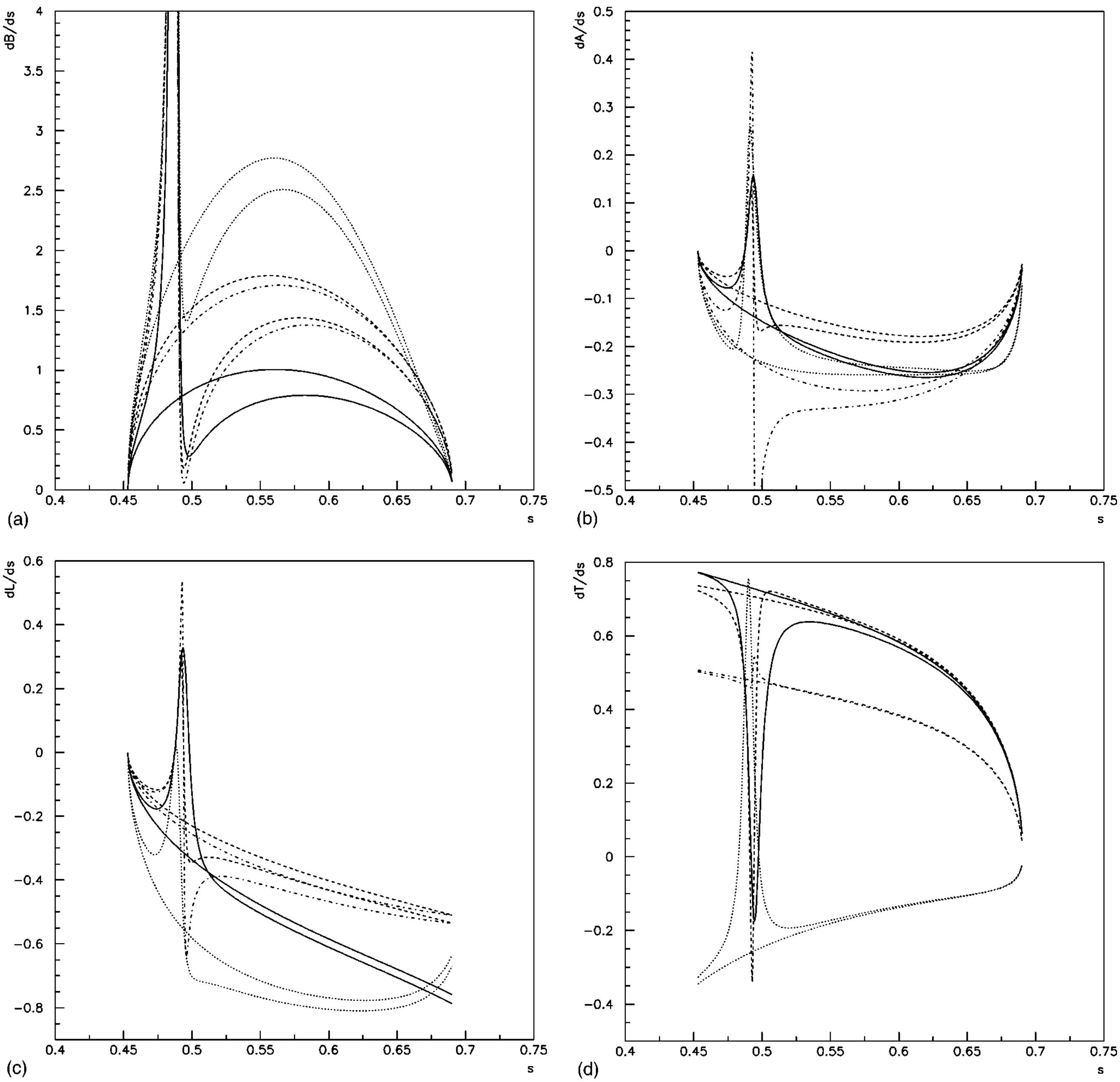

FIG. 4. The IMS(a), FBA(b), P_L(c), and P_T(d) of the process $B\rightarrow K^*\tau^+\tau^-$. The line conventions are the same as given in the legend of Fig. 1.

butions of NHB's.

Normal polarizations for both $B\rightarrow Kl^+l^-$ and $B\rightarrow K^*l^+l^-$ are small and can be neglected because the imaginary parts of Wilson coefficients are small in SUSY models without CP-violating phases which are implicitly assumed in the paper.

The behavior of IMS (a), FBA (b), P_L (c), and P_T (d) shown Figs. 1–4 can be understood with the formula given in Sec. III. With Eqs. (3.55), (3.10), and (3.11), we see that the contributions of NHB's are contained in the terms of $\mathcal{S}_1$ and D'. At the high $\hat{s}$ regions, it is these two terms which are important. This explained the behavior of IMS given in (a) of Figs. 1 and 2. Equation (3.58) shows that the FBA is proportional to the mass of the lepton. For the case $B\rightarrow K\mu^+\mu^-$, due to smallness of the mass μ, the FBA does not vanish but it is hard to measure. While for the case $B\rightarrow K\tau^+\tau^-$, the mass τ is quite large and observing FBA is relatively easy. For SUSY II, though the numerator of FBA is comparatively large, the large IMS suppresses the value of FBA; for SUSY III, the numerator is relatively small, but the FBA's do demonstrate the effects of NHB's more manifestly, as shown in Fig. 2(b) due to the smallness of IMS. Equations (3.63) and (3.64) show that for the case $l=\mu$, the contributions of NHB's to P_N, P_T are suppressed by the mass of μ. But for the case $l=\tau$, the contributions of NHB's become quite manifest both for SUSY II and SUSY III. The term with D' in Eq. (3.63) will change its sign when there exist relatively not too small contributions of NHB's, the fact deduced from

TABLE V. Partial decay widths for $B \to K\mu^+\mu^-$. LCSR means the approach light-cone QCD sum rules, SVZ means the SVZ QCD sum rule [17]. Character A means the region $[\hat{s}_0,(\hat{m}_\psi-\hat{\delta})^2]$, B $[(\hat{m}_\psi-\hat{\delta})^2,(\hat{m}_\psi+\hat{\delta})^2]$, C $[(\hat{m}_\psi+\hat{\delta})^2,(\hat{m}_{\psi'}-\hat{\delta})^2]$, D $[(\hat{m}_{\psi'}-\hat{\delta})^2,(\hat{m}_{\psi'}+\hat{\delta})^2]$, and E $[(\hat{m}_{\psi'}+\hat{\delta})^2,\hat{s}^2_{\max}]$. The unit is $\Gamma_B\times 10^{-6}$, which is 4.22×10^{-19} GeV. δ is selected to be 0.2 GeV. $\hat{\delta}$ is normalized with M_B

Model		A	B	C	D	E	tot(SD)	tot(SD+LD)
SM	LCSR	0.353	54.707	0.032	4.566	0.076	0.573	59.736
	SVZ	0.215	22.918	0.015	1.593	0.026	0.299	24.767
SUSY I	LCSR	0.425	54.723	0.037	4.576	0.086	0.675	59.847
	SVZ	0.179	22.910	0.011	1.586	0.019	0.236	24.704
SUSY II	LCSR	0.556	54.865	0.131	4.833	0.849	2.067	61.233
	SVZ	0.348	23.009	0.068	1.726	0.321	1.002	25.473
SUSY III	LCSR	0.429	54.727	0.040	4.584	0.109	0.717	59.889
	SVZ	0.181	22.912	0.012	1.590	0.028	0.255	24.723

Eq. (3.10), that explains why the sign of P_T is changed. The difference between cases SUSY II and SUSY III is small, the reason is just the same as stated in the analysis of FBA.

Since the terms incorporating the contributions of NHB's are proportional to λ as shown in Eq. (3.57), which approaches zero at high $\hat{s}$ regions; while at small $\hat{s}$ regions, the effects of NHB's are dwarfed by the other contributions. Therefore, only when C_{Q_i} are quite large could the effects of NHB's be manifest, as shown in Figs. 3(a) and 4(a). According to Eq. (3.59), at high $\hat{s}$ regions, the effects of NHB's would be suppressed by λ and $1-\hat{s}-\hat{m}^2_{K^*}$. The same suppression mechanism exists for P_L. This suppression mechanism explains the fact that the processes $B \to K^* l^+ l^-$ are not sensitive to the effects of NHB's. However, when there exist large contributions of NHB's, the sign of P_T will be changed, as indicated in both Figs. 3(d) and 4(d).

The partial decay widths (PDW's) are listed in Tables V–VIII. We see that at the high $\hat{s}$ region, for the process $B \to K l^+ l^-$, $l=\mu,\tau$, the contributions of NHB's do show up, as expected. For $B \to K^* l^+ l^-$, the effects of NHB's in the high $\hat{s}$ region is significant when $l=\tau$, while they are small for $l=\mu$. It can be read out from these four tables that the results are consistent with Figs. 1(a), 2(a), 3(a), and 4(a). In order to estimate the theoretical uncertainty brought by the methods calculating the weak form factors, we use the form factors calculated with LCSR and SVZ QCD sum rules (SVZ) method [17]. For $B \to K l^+ l^-$, PDW's calculated with form factors obtained by the SVZ method are 50% of those by the LCSR approach, while for $B \to K^* l^+ l^-$, PDW's increase 100% or so. We see that at low $\hat{s}$ regions the theoretical uncertainty can reach from 100% to 200%. Another point worthy of mention is that the contribution of resonances dominate the integerated decay width, as had been pointed out in [29].

V. CONCLUSION

We have calculated invariant mass spectrum, FBA's, and lepton polarizations for $B \to K l^+ l^-$ and $B \to K^* l^+ l^-$ $l=\mu,\tau$ in SUSY theories. In particular, we have analyzed the effects of NHB's on these processes. It is shown that the effects of the NHB's on $B \to K\tau^+\tau^-$ and $B \to K^*\tau^+\tau^-$ in some regions of the parameter space of SUSY models are considerable and remarkable. The reason lies in the mass of the τ, which can magnify the effects of NHB's and can be seen from the related formula. The numerical results imply that there still exist possiblities to observe the effects of NHB in $B \to K\mu^+\mu^-$ and $B \to K^*\mu^+\mu^-$ through IMS, FBA, and lepton polarizations of these processes, in particular, for $B \to K\mu^+\mu^-$ in the case of SUSY II. The partial width in the high $\hat{s}$ where short distance physics dominates can be enhanced by a factor of 12 compared to SM. Our analysis also show that the theoretical uncertainties brought in calculating the weak form factors are quite large. But the effects of NHB's will not be washed out and can stand out in

TABLE VI. Partial decay widths for $B \to K^*\mu^+\mu^-$. Other conventions can be found in Table V.

Model		A	B	C	D	E	tot(SD)	tot(SD+LD)
SM	LCSR	0.930	83.257	0.141	9.976	0.258	1.882	94.562
	SVZ	2.943	111.278	0.147	7.504	0.137	3.639	122.008
SUSY I	LCSR	1.627	83.402	0.198	10.085	0.330	2.915	95.64
	SVZ	4.517	111.423	0.183	7.552	0.149	5.291	123.825
SUSY II	LCSR	1.178	83.431	0.234	10.164	0.352	2.677	95.360
	SVZ	2.801	111.292	0.156	7.525	0.145	3.522	121.918
SUSY III	LCSR	1.631	83.407	0.201	10.092	0.334	2.938	95.664
	SVZ	4.518	111.425	0.184	7.553	0.150	5.296	123.830

TABLE VII. Partial decay widths of $B\to K\tau^+\tau^-$. A' means $[\hat{s}_0,(\hat{m}_{\psi}-\hat{\delta})^2]$, B' means $[(\hat{m}_{\psi'}+\hat{\delta})^2,\hat{s}_{\max}]$. The unit is $\Gamma_B \times 10^{-6}$, which is 4.22×10^{-19} GeV.

Model		A'	B'	tot(SD)	tot(SD+LD)
SM	LCSR	1.884	0.094	0.132	1.978
	SVZ	0.659	0.036	0.054	0.695
SUSY I	LCSR	1.884	0.086	0.131	1.970
	SVZ	0.655	0.025	0.038	0.680
SUSY II	LCSR	2.022	1.496	1.674	3.519
	SVZ	0.726	0.552	0.637	1.278
SUSY III	LCSR	1.874	0.094	0.129	1.968
	SVZ	0.651	0.026	0.035	0.677

TABLE VIII. Partial decay widths of $B\to K^*\tau^+\tau^-$. Other conventions can be found in Table VII.

Model		A'	B'	tot(SD)	tot(SD+LD)
SM	LCSR	4.045	0.096	0.183	4.141
	SVZ	3.029	0.048	0.102	3.076
SUSY I	LCSR	4.088	0.173	0.327	4.261
	SVZ	3.052	0.072	0.159	3.124
SUSY II	LCSR	4.148	0.266	0.460	4.413
	SVZ	3.054	0.084	0.167	3.138
SUSY III	LCSR	4.078	0.168	0.312	4.246
	SVZ	3.050	0.071	0.156	3.121

some regions of the parameter space in MSSM. If only partial widths are measured, it is difficult to observe the effects of NHB's except for the decay $B\to K\tau^+\tau^-$. However, the combined analysis of IBS, FBA, and lepton polarizations can provide useful knowledge when looking for SUSY. Finally, we would like to point out that FBA for $B\to Kl^+l^-$ vanishes (or, more precisely, is negligibly small) in SM and it does not vanish in 2HDM and SUSY models with large $\tan\beta$ due to the contributions of NHB's. However, only in SUSY models and for $l=\tau$ is it large enough to be observed in B factories in the near future.

Note added in proof. Before this work was finished, we noticed in paper [43], which points out that in 2HDM, the missed box diagram can preserve the gauge invariance of the effective Hamiltonian and can contribute considerably. However, in SUSY, the dominant contribution is from the SUSY self-energy diagram which is proportional to $\tan^3\beta$; therefore, although they are important to a certain concern, numerically the missed boxed diagram will not change our conclusion about the contribution of NHB's in SUSY.

ACKNOWLEDGMENTS

Q.S.Y. would like to thank Dr. W.J. Huo for his help during the work. This work was supported in part by the National Nature Science Foundation of China.

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

[2] T. Goto *et al.*, Phys. Lett. B **460**, 333 (1999); S. Baek and P. Ko, *ibid.* **462**, 95 (1999); Y.G. Kim, P. Ko, and J.S. Lee, Nucl. Phys. **B544**, 64 (1999), and references therein. For the earlier references, see, for example, the references in Ref. [9].

[3] Y.B. Dai, C.S. Huang, and H.W. Huang, Phys. Lett. B **390**, 257 (1997).

[4] C.S. Huang and Q.S. Yan, Phys. Lett. B **442**, 209 (1998); C.S. Huang, W. Liao, and Q.S. Yan, Phys. Rev. D **59**, 011701 (1999).

[5] S. Fukae, C.S. Kim, and T. Yoshikawa, Phys. Rev. D **61**, 074015 (2000).

[6] S. Fukae, C.S. Kim, T. Morozumi, and T. Yoshikawa, Phys. Rev. D **59**, 074013 (1999).

[7] Y.G. Kim, P. Ko, and J.S. Lee, Nucl. Phys. **B544**, 64 (1999).

[8] T. Goto *et al.*, Phys. Rev. D **55**, 4273 (1997); T. Goto, Y. Okada, and Y. Shimizu, *ibid.* **58**, 094006 (1998).

[9] E. Lunghi *et al.*, Nucl. Phys. **B568**, 120 (2000).

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

[11] Y. Grossman, Z. Ligeti, and E. Nardi, Phys. Rev. D **55**, 2768 (1997).

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

[13] S. Choudhury *et al.*, Phys. Rev. D **60**, 115004 (1999).

[14] Chao-Shang Huang and Shou-Hua Zhu, Phys. Rev. D **61**, 015011 (2000).

[15] I.I. Bigi, M. Shifman, N.G. Vraltsev, and A.I. Vainstein, Phys. Rev. Lett. **71**, 496 (1993); B. Blok, L. Kozrakh, M. Shifman, and A.I. Vainstein, Phys. Rev. D **49**, 3356 (1994); A.V. Manohar and M.B. Wise, *ibid.* **49**, 1310 (1994); S. Balk, T.G. Körner, D. Pirjol, and K. Schilcher, Z. Phys. C **64**, 37 (1994); A.F. Falk, Z. Ligeti, M. Neubert, and Y. Nir, Phys. Lett. B **326**, 145 (1994).

[16] W. Jaus and D. Wyler, Phys. Rev. D **41**, 3405 (1990); P. Colangelo *et al.*, Phys. Lett. B **317**, 183 (1993); Ceng *et al.*, Phys. Rev. D **54**, 3656 (1996).

[17] P. Colangelo *et al.*, Phys. Rev. D **53**, 3672 (1996); **57**, 3186(E) (1998).

[18] I.I. Balitsky, V.M. Braun, and A.V. Kolesnichenko, Nucl. Phys. **B312**, 509 (1989).

[19] V.L. Chernyak and I.R. Zhitnitsky, Nucl. Phys. **B345**, 137 (1990).

[20] V.M. Braun, Preprint NORDITA–98–1–P, hep-ph/9801222.

[21] A. Khodjamirian and R. Rückl, Preprint WUE–ITP–97–049, hep-ph/9801443.

[22] P. Ball and V.M. Braun, Phys. Rev. D **58**, 094016 (1998); Nucl. Phys. **B543**, 201 (1999); P. Ball *et al.*, *ibid.* **B529**, 323 (1998); P. Ball, J. High Energy Phys. **09**, 005 (1998); **01**, 010 (1999).

[23] A. Ali, P. Ball, L.T. Handoko, and G. Hiller, Phys. Rev. D **61**, 074024 (2000).

[24] CLEO Collaboration, S. Glenn *et al.*, Phys. Rev. D **60**, 034017 (1999); CLEO Collaboration, T. Affolder *et al.*, Phys. Rev. Lett. **83**, 3378 (1999).

[25] F. Krüger and J. C. Romão, Phys. Rev. D **62**, 034020 (2000).
[26] N.G. Deshpande and J. Trampetić, Phys. Rev. Lett. **60**, 2583 (1988).
[27] G. Burdman, Phys. Rev. D **52**, 6400 (1995); W. Roberts, *ibid.* **54**, 863 (1996).
[28] D.S. Du and C. Liu, Phys. Lett. B **317**, 179 (1993).
[29] D.S. Liu, Phys. Rev. D **52**, 5056 (1995); M.R. Ahmady, D.S. Liu, and A.H. Fariborz, Z. Phys. C **71**, 301 (1996); M.R. Ahmady and D.S. Liu, Phys. Lett. B **324**, 231 (1994).
[30] T.M. Aliev, C.S. Kim, and Y.G. Kim, Phys. Rev. D **62**, 014026 (2000).
[31] M. Jezabek and P. Urban, Nucl. Phys. **B525**, 350 (1998).
[32] D. Melikhov, N. Nikitin, and S. Simula, Phys. Lett. B **442**, 381 (1998).
[33] C. Greub, A. Ioannissian, and D. Wyler, Phys. Lett. B **346**, 149 (1995).
[34] T.M. Aliev and M. Savci, Phys. Lett. B **481**, 275 (2000).
[35] T. M. Aliev, M. Savci, A. Özpineci, and H. Koru, J. Phys. G **24**, 49 (1998).
[36] C.Q. Geng and C.P Kao, Phys. Rev. D **54**, 5636 (1996).
[37] C.S. Huang *et al.*, hep-ph/9810412; Commun. Theor. Phys. **32**, 499 (1999).
[38] S. Bertolini, F. Borzynatu, A. Masiero, and G. Ridolfi, Nucl. Phys. **B353**, 591 (1991).
[39] G. Buchalla, A.J. Buras, and M.E. Lauthenbache, Rev. Mod. Phys. **68**, 1125 (1996); A. Buras *et al.*, Nucl. Phys. **B400**, 37 (1993); **B400**, 75 (1993); S. Bertonlini *et al.*, *ibid.* **B353**, 591 (1991).
[40] A. Ali, G.F. Giudice, and T. Mannel, Z. Phys. C **67**, 417 (1995).
[41] A.J. Buras and M. Münz, Phys. Rev. D **52**, 186 (1995).
[42] A. Ali, T. Mannel, and T. Morozumi, Phys. Lett. B **273**, 505 (1991).
[43] H.E. Logan and U. Nierste, hep-ph/0004139.