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Hadron spectra and spectral moments in the decay $B\to X_s l^+ l^-$ using HQET

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We compute the leading order (in α_s) perturbative QCD and power ($1/m_b^2$) corrections to the hadronic invariant mass and hadron energy spectra in the decay $B\to X_s l^+ l^-$ in standard model using the heavy quark expansion technique (HQET). Results for the first two hadronic moments $\langle S_H^n\rangle$ and $\langle E_H^n\rangle$, $n=1,2$, are presented here working out their sensitivity on the HQET parameters λ_1 and $\bar{\Lambda}$. Data from the forthcoming B facilities can be used to measure the short-distance contribution in $B\to X_s l^+ l^-$ and determine the HQET parameters from the moments $\langle S_H^n\rangle$. This can be combined with the analysis of semileptonic decays $B\to X l \nu_l$ to determine them precisely. [S0556-2821(98)50417-1]

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The semileptonic inclusive decays $B\to X_s l^+ l^-$, where $l^\pm = e^\pm, \mu^\pm, \tau^\pm$, offer, together with the radiative electromagnetic penguin decay $B\to X_s+\gamma$, presently the most popular testing grounds for the standard model (SM) in the flavor sector [1]. In the quest of developing a precise theory for flavor changing weak interactions, we study the hadron spectra and hadron spectral moments in $B\to X_s l^+ l^-$ using the heavy quark expansion technique HQET [2–4]. The study of the decay $B\to X l \nu_l$ in this context has received a lot of theoretical attention [5–10].

We have computed the leading order (in α_s) perturbative QCD and power ($1/m_b^2$) corrections to the hadronic invariant mass and hadron energy spectra in the decay $B\to X_s l^+ l^-$. Relegating the details of the derivation to a companion publication [11], here the power- and perturbatively corrected hadronic spectral moments $\langle S_H^n\rangle$ and $\langle E_H^n\rangle$ are presented for $n=1,2$. The former are sensitive to the HQET parameters $\bar{\Lambda}$ and λ_1 and we work out this dependence numerically, showing that these moments would provide an independent determination of the HQET parameters in $B\to X_s l^+ l^-$. We argue that a simultaneous analysis of the moments and spectra in $B\to X_s l^+ l^-$ and $B\to X l \nu_l$ will allow us to determine the HQET parameters with a high precision.

We start with the definition of the kinematics of the decay at parton level, $b(p_b)\to s(p_s)(+g(p_g))+l^+(p_+)+l^-(p_-)$, where g denotes a gluon from the $O(\alpha_s)$ correction. The corresponding kinematics at the hadron level is written as: $B(p_B)\to X_s(p_H)+l^+(p_+)+l^-(p_-)$. We define by q the momentum transfer to the lepton pair $q=p_+ + p_-$ and $s\equiv q^2$; the 4-vector v denotes the velocity of both the b-quark and the B-meson, $p_b=m_b v$ and $p_B=m_B v$ and further $u=-(p_b-p_+)^2+(p_b-p_-)^2$. The hadronic invariant mass and the hadron energy in the final state is denoted by S_H and E_H, respectively; corresponding quantities at parton level are the invariant mass s_0 and the scaled parton energy $x_0\equiv E_0/m_b$. From energy-momentum conservation, the following equalities hold in the b-quark, equivalently B-meson, rest frame [$v=(1,0,0,0)$]:

$$x_0 = 1 - v\cdot\hat{q}, \quad \hat{s}_0 = 1 - 2v\cdot\hat{q} + \hat{s},$$
$$E_H = m_B - v\cdot q, \quad S_H = m_B^2 - 2m_B v\cdot q + s, \tag{1}$$

where dimensionless variables with a hat are scaled by the b-quark mass, e.g., $\hat{s}=s/m_b^2$, $\hat{m}_s = m_s/m_b$, etc. The relation between the B meson and b quark mass is given by the HQET mass relation $m_B = m_b + \bar{\Lambda} - 1/2m_b(\lambda_1 + 3\lambda_2) + \cdots$, where the ellipses denote terms higher order in $1/m_b$. The quantity λ_2 is known precisely from the $B^* - B$ mass difference, with $\lambda_2 \simeq 0.12$ GeV^2 [2–4]. A precise determination of the other two parameters is of interest here.

The effective Hamiltonian governing the decay $B\to X_s l^+ l^-$ is given as [12]

$$\mathcal{H}_{\rm eff}(b\to s+X, X=\gamma, l^+l^-) = -\frac{4G_F}{\sqrt{2}} V_{ts}^* V_{tb}\left[\sum_{i=1}^{6} C_i(\mu) O_i + C_7(\mu)\frac{e}{16\pi^2}\bar{s}_\alpha \sigma_{\mu\nu}(m_b R + m_s L) b_\alpha F^{\mu\nu} + C_8(\mu) O_8 + C_9(\mu)\frac{e^2}{16\pi^2}\bar{s}_\alpha \gamma^\mu L b_\alpha \bar{l}\gamma_\mu l + C_{10}\frac{e^2}{16\pi^2}\bar{s}_\alpha \gamma^\mu L b_\alpha \bar{l}\gamma_\mu \gamma_5 l\right], \tag{2}$$

where $L(R)=1/2(1\mp\gamma_5)$, V_{ij} are the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements. Note that the chromomagnetic operator O_8 does not contribute to the decay $B\to X_s l^+ l^-$ in the approximation which we use here. The values of the Wilson

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coefficients used in the numerical calculations are $C_1=-0.24$, $C_2=1.103$, $C_3=0.011$, $C_4=-0.025$, $C_5=0.007$, $C_6=-0.03$, $C_7^{\rm eff}=-0.311$, $C_9=4.153$, $C_{10}=-4.546$. Here, $C_7^{\rm eff}\equiv C_7-C_5/3-C_6$, and for C_9 we use the naive dimensional regularization (NDR) scheme.

One can express the Dalitz distribution in $B\to X_s l^+ l^-$ as

$$\frac{d\Gamma}{d\hat{u}d\hat{s}d(v\cdot\hat{q})}=\frac{1}{2m_B}\frac{{G_F}^2\alpha^2}{2\pi^2}\frac{m_b^4}{256\pi^4}|V_{ts}^*V_{tb}|^2 2\,\mathrm{Im}(T^L{}_{\mu\nu}L^{L\mu\nu}+T^R{}_{\mu\nu}L^{R\mu\nu}), \tag{3}$$

and the hadronic and leptonic tensors $T^{L/R}{}_{\mu\nu}$ and $L^{L/R\mu\nu}$ are given in [12]. Using Lorentz decomposition, the tensor $T_{\mu\nu}$ can be expanded in terms of three structure functions T_i,

$$T_{\mu\nu}=-T_1 g_{\mu\nu}+T_2 v_\mu v_\nu+T_3 i\varepsilon_{\mu\nu\alpha\beta}v^\alpha\hat{q}^\beta, \tag{4}$$

and the ones which do not contribute to the amplitude for massless leptons have been neglected.

Concerning the $O(\alpha_s)$ corrections, note that only the matrix element of the operator $O_9\equiv e^2/(16\pi^2)\bar{s}_\alpha\gamma^\mu L b_\alpha \bar{l}\gamma_\mu l$ is subject to such corrections. These can be obtained by using the existing results in the literature by decomposing the vector current in O_9 as $V=(V-A)/2+(V+A)/2$. The $(V-A)$ and $(V+A)$ currents yield the same hadron energy spectrum [13] and there is no interference term present in this spectrum for massless leptons. So, the correction for the vector current case can be taken from the corresponding result for the charged $(V-A)$ case [14,15].

We have calculated the order α_s perturbative QCD correction for the hadronic invariant mass in the range $\hat{m}_s^2<\hat{s}_0\leqslant 1$. Since the decay $b\to s+l^++l^-$ contributes in the parton model only at $\hat{s}_0=\hat{m}_s^2$, only the bremsstrahlung graphs $b\to s+g+l^++l^-$ contribute in this range. Also for this distribution, the results can be taken from the existing literature. We use the Sudakov exponentiated double differential decay rate $d^2\Gamma/dxdy$, derived for the decay $B\to X_u l\nu_l$ in [7], which we have checked, after changing the normalization for $B\to X_s l^+l^-$. The most significant effect of the bound state is in the difference between m_B and m_b which is dominated by $\bar{\Lambda}$. The spectrum $d\mathcal{B}/dS_H$ obtained is valid in the region $m_B(m_B\bar{\Lambda}-\bar{\Lambda}^2+m_s^2)/(m_B-\bar{\Lambda})\leqslant S_H\leqslant m_B^2$ (or $m_B\bar{\Lambda}\leqslant S_H\leqslant m_B^2$, neglecting m_s) which excludes the zeroth order and virtual gluon kinematics $(s_0=m_s^2)$. The hadronic invariant mass spectrum thus found depends rather sensitively on m_b (or equivalently $\bar{\Lambda}$). An analogous analysis for the decay $B\to X_\mu l\nu_l$ has been performed in [8].

The hadronic tensor in Eq. (4) can be expanded in inverse powers of m_b with the help of the HQET techniques [2–4]. The leading term in this expansion, i.e., $\mathcal{O}(m_b^0)$, reproduces the parton model result [16,17]. In HQET, the next to leading power corrections are parametrized in terms of λ_1 and λ_2. The contributions of the power corrections to the structure functions T_i have been calculated up to $O(1/m_b^3)$ in [12]. After contracting the hadronic and leptonic tensors and with the help of the kinematic identities given in Eq. (1), we can make the dependence on x_0 and $\hat{s}_0$ explicit,

$$T^{L/R}{}_{\mu\nu}L^{L/R\mu\nu}=m_b^2\bigg\{2(1-2x_0+\hat{s}_0)T_1^{L/R}+\left[x_0^2-\frac{1}{4}\hat{u}^2-\hat{s}_0\right]T_2^{L/R}\mp(1-2x_0+\hat{s}_0)\hat{u}T_3^{L/R}\bigg\},$$

and derive the double differential power corrected spectrum by integrating Eq. (3) over $\hat{u}$ [11]:

$$\frac{d^2\mathcal{B}}{dx_0 d\hat{s}_0}=-\frac{8}{\pi}\mathcal{B}_0\mathrm{Im}\sqrt{x_0^2-\hat{s}_0}\bigg\{(1-2x_0+\hat{s}_0)T_1(\hat{s}_0,x_0)+\frac{x_0^2-\hat{s}_0}{3}T_2(\hat{s}_0,x_0)\bigg\}+\mathcal{O}(\lambda_i\alpha_s). \tag{5}$$

As the structure function T_3 does not contribute to the branching ratio, we did not consider it in our present work. The functions $T_1(\hat{s}_0,x_0)$ and $T_2(\hat{s}_0,x_0)$ can be seen together with other details of the calculations in [11].

The branching ratio for $B\to X_s l^+l^-$ is usually expressed in terms of the measured semileptonic branching ratio $\mathcal{B}_{sl}$ for the decay $B\to X_c l v_l$. This fixes the normalization constant $\mathcal{B}_0$ to be

$$\mathcal{B}_0\equiv\mathcal{B}_{sl}\frac{3\alpha^2}{16\pi^2}\frac{|V_{ts}^*V_{tb}|^2}{|V_{cb}|^2}\frac{1}{f(\hat{m}_c)\kappa(\hat{m}_c)}, \tag{6}$$

where $f(\hat{m}_c)$ is the phase space factor for $\Gamma(B\to X_c l\nu_l)$ and $\kappa(\hat{m}_c)$ accounts for both the $O(\alpha_s)$ QCD correction to the semileptonic decay width [18] and the leading order $(1/m_b)^2$ power correction [2]. They are given explicitly in [5]. The hadron energy spectrum can now be obtained by integrating over $\hat{s}_0$ with the kinematic boundaries: $\max(\hat{m}_s^2,-1+2x_0+4\hat{m}_l^2)\leqslant\hat{s}_0\leqslant x_0^2$, $\hat{m}_s\leqslant x_0\leqslant\frac{1}{2}(1+\hat{m}_s^2-4\hat{m}_l^2)$. Here we keep $\hat{m}_l$ as a regulator wherever it is necessary. Including the leading power corrections, the hadron energy spectrum in the decay $B\to X_s l^+l^-$ is derived by us and given in [11].

The lowest spectral moments in the decay $B\to X_s l^+l^-$ at the parton level are worked out by taking into account the two types of corrections discussed earlier, namely the leading power $1/m_b$ and the perturbative $\mathcal{O}(\alpha_s)$ corrections. To that end, we define for integers n and m:

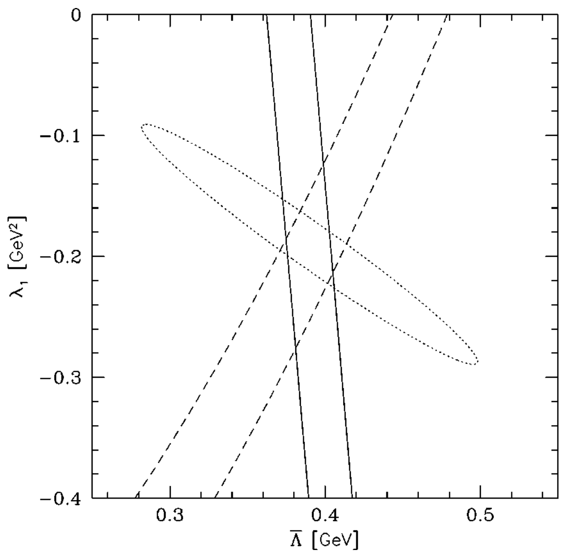

FIG. 1. $\langle S_H\rangle$ (solid bands) and $\langle S_H^2\rangle$ (dashed bands) correlation in $(\lambda_1-\bar{\Lambda})$ space corresponding to the $(\mu^+\mu^-)$ values in Table I. the correlation from $B\to Xl\nu_l$ [6] is also shown here (ellipse).

$$\mathcal{M}_{l^+l^-}^{(n,m)}\equiv\frac{1}{\mathcal{B}_0}\int(\hat{s}_0-\hat{m}_s^2)^n x_0^m\frac{d\mathcal{B}}{d\hat{s}_0\,dx_0}\,d\hat{s}_0 dx_0, \qquad (7)$$

which obey $\langle x_0^m(\hat{s}_0-\hat{m}_s^2)^n\rangle=(\mathcal{B}_0/\mathcal{B})\mathcal{M}_{l^+l^-}^{(n,m)}$. They can be expanded in α_s and $1/m_b$:

$$\mathcal{M}_{l^+l^-}^{(n,m)}=D_0^{(n,m)}+\frac{\alpha_s}{\pi}C_9^2A^{(n,m)}+\hat{\lambda}_1D_1^{(n,m)}+\hat{\lambda}_2D_2^{(n,m)}, \qquad (8)$$

with a further decomposition into pieces from different Wilson coefficients for $i=0,1,2$:

$$D_i^{(n,m)}=\alpha_i^{(n,m)}C_7^{\mathrm{eff}^2}+\beta_i^{(n,m)}C_{10}^2+\gamma_i^{(n,m)}C_7^{\mathrm{eff}}+\delta_i^{(n,m)}. \qquad (9)$$

The terms $\gamma_i^{(n,m)}$ and $\delta_i^{(n,m)}$ in Eq. (9) result from the terms proportional to $\mathrm{Re}(C_9^{\mathrm{eff}})C_7^{\mathrm{eff}}$ and $|C_9^{\mathrm{eff}}|^2$ in Eq. (5), respectively. The explicit expressions for $\alpha_i^{(n,m)}$, $\beta_i^{(n,m)}$, $\gamma_i^{(n,m)}$, $\delta_i^{(n,m)}$ are given in [11].

TABLE I. Hadronic spectral moments for $B\to X_s\mu^+\mu^-$ and $B\to X_se^+e^-$ in HQET with $\bar{\Lambda}=0.39$ GeV and $\lambda_1=-0.2$ GeV^2. The errors result by varying μ, α_s and the m_t within their stated ranges in text.

HQET	$\langle S_H\rangle$ (GeV^2)	$\langle S_H^2\rangle$ (GeV^4)	$\langle E_H\rangle$ (GeV)	$\langle E_H^2\rangle$ (GeV^2)
$\mu^+\mu^-$	1.64 ± 0.06	4.48 ± 0.29	2.21 ± 0.04	5.14 ± 0.16
e^+e^-	1.79 ± 0.07	4.98 ± 0.29	2.41 ± 0.06	6.09 ± 0.29

The leading perturbative contributions for the hadronic invariant mass and hadron energy moments can be obtained analytically,

$$A^{(0,0)}=\frac{25-4\pi^2}{9},\quad A^{(1,0)}=\frac{91}{675},\quad A^{(2,0)}=\frac{5}{486},$$
$$A^{(0,1)}=\frac{1381-210\pi^2}{1350},\quad A^{(0,2)}=\frac{2257-320\pi^2}{5400}. \qquad (10)$$

The zeroth moment $n=m=0$ is needed for the normalization; the result for $A^{(0,0)}$ was first derived in [18]. Likewise, the first mixed moment $A^{(1,1)}$ can be extracted from results for the decay $B\to Xl\nu_l$ [5] after changing the normalization, $A^{(1,1)}=3/50$. For the lowest order parton model contribution $D_0^{(n,m)}$, we find, in agreement with [5], that the first two hadronic invariant mass moments $\langle\hat{s}_0-\hat{m}_s^2\rangle$, $\langle(\hat{s}_0-\hat{m}_s^2)^2\rangle$ and the first mixed moment $\langle x_0(\hat{s}_0-\hat{m}_s^2)\rangle$ vanish: $D_0^{(n,0)}=0$, for $n=1,2$ and $D_0^{(1,1)}=0$.

Using the expressions for the HQET moments derived by us [11], we present the numerical results for the hadronic moments in $B\to X_sl^+l^-$. The parameters used are $m_s=0.2$ GeV, $m_c=1.4$ GeV, $m_b=4.8$ GeV, $m_t=175\pm5$ GeV, $\mu=m_b{}^{+m_b}_{-m_b/2}$, $\alpha_s(m_Z)=0.117\pm0.005$, $\alpha^{-1}=129$. We find for the short-distance hadronic moments, valid up to $\mathcal{O}(\alpha_s/m_B^2,1/m_B^3)$:

$$\langle S_H\rangle=m_B^2\left(\frac{m_s^2}{m_B^2}+0.093\frac{\alpha_s}{\pi}-0.069\frac{\bar{\Lambda}}{m_B}\frac{\alpha_s}{\pi}+0.735\frac{\bar{\Lambda}}{m_B}+0.243\frac{\bar{\Lambda}^2}{m_B^2}+0.273\frac{\lambda_1}{m_B^2}-0.513\frac{\lambda_2}{m_B^2}\right),$$
$$\langle S_H^2\rangle=m_B^4\left(0.0071\frac{\alpha_s}{\pi}+0.138\frac{\bar{\Lambda}}{m_B}\frac{\alpha_s}{\pi}+0.587\frac{\bar{\Lambda}^2}{m_B^2}-0.196\frac{\lambda_1}{m_B^2}\right),$$
$$\langle E_H\rangle=0.367m_B\left(1+0.148\frac{\alpha_s}{\pi}-0.352\frac{\bar{\Lambda}}{m_B}\frac{\alpha_s}{\pi}+1.691\frac{\bar{\Lambda}}{m_B}+0.012\frac{\bar{\Lambda}^2}{m_B^2}+0.024\frac{\lambda_1}{m_B^2}+1.070\frac{\lambda_2}{m_B^2}\right),$$
$$\langle E_H^2\rangle=0.147m_B^2\left(1+0.324\frac{\alpha_s}{\pi}-0.128\frac{\bar{\Lambda}}{m_B}\frac{\alpha_s}{\pi}+2.954\frac{\bar{\Lambda}}{m_B}+2.740\frac{\bar{\Lambda}^2}{m_B^2}-0.299\frac{\lambda_1}{m_B^2}+0.162\frac{\lambda_2}{m_B^2}\right), \qquad (11)$$

where the numbers shown correspond to central values of the parameters. Further, using $\alpha_s(m_b)=0.21$ and $\lambda_2=0.12\ \mathrm{GeV}^2$, the explicit dependence of the hadronic moments given in Eq. (11) on the HQET parameters λ_1 and $\bar{\Lambda}$ can be worked out:

$$\begin{aligned}\langle S_H\rangle &= 0.0055 m_B^2\left(1+132.61\frac{\bar{\Lambda}}{m_B}+44.14\frac{\bar{\Lambda}^2}{m_B^2}+49.66\frac{\lambda_1}{m_B^2}\right),\\ \langle S_H^2\rangle &= 0.00048 m_B^4\left(1+19.41\frac{\bar{\Lambda}}{m_B}+1223.41\frac{\bar{\Lambda}^2}{m_B^2}-408.39\frac{\lambda_1}{m_B^2}\right).\end{aligned} \tag{12}$$

The dependence of the energy moments $\langle E_H^n\rangle$ on $\bar{\Lambda}$ and λ_1 is very weak and we do not show these here. We have calculated numerically the hadronic moments in HQET for the decay $B\to X_s l^+ l^-$, $l=\mu,e$. The theoretical errors on these moments following from the errors on the input parameters m_t, α_s and the scale μ were estimated by varying these parameters in the indicated $\pm 1\sigma$ ranges, one at a time, and adding the individual errors in quadrature. They are presented in Table I where we have used $\bar{\Lambda}=0.39$ GeV and $\lambda_1=-0.2\ \mathrm{GeV}^2$. The correlations on the HQET parameters λ_1 and $\bar{\Lambda}$ which follow from (assumed) fixed values of the hadronic invariant mass moments $\langle S_H\rangle$ and $\langle S_H^2\rangle$ are shown in Fig. 1 (for the decay $B\to X_s\mu^+\mu^-$). As the entries in Table I are calculated for the best-fit values of λ_1 and $\bar{\Lambda}$ taken from the analysis of Gremm *et al.* [6] for the electron energy spectrum in $B\to X l \nu_l$, which is shown as an ellipse in Fig. 1, there is no surprise that these curves meet at this point. It is, however, clear that the constraints from the decays $B\to X_s l^+ l^-$ and $B\to X l \nu_l$ are complementary. Using the CLEO cuts on hadronic and dileptonic masses [19], we estimate that $O(200)$ $B\to X_s l^+ l^-(l=e,\mu)$ events will be available per 10^7 $B\bar{B}$ hadrons [11]. So, there will be plenty of $B\to X_s l^+ l^-$ decays in the forthcoming B facilities to measure the correlation shown in Fig. 1.

The theoretical stability of the moments has to be checked against higher order corrections and the error estimates presented here will have to be improved. The ''BLM-enhanced'' two-loop corrections [20] can be included at the parton level [5,21], but not being crucial to our point we have not done this. More importantly, corrections in $1/m_b^3$ are not included here. The second moment $\langle S_H^2\rangle$ is susceptible to the presence of $1/m_b^3$ corrections as shown for the decay $B\to X l \nu_l$ [9]. This will considerably enlarge the theoretical error represented by the dashed band for $\langle S_H^2\rangle$ in Fig. 1. Fortunately, the coefficient of the $\bar{\Lambda}/m_B$ term in $\langle S_H\rangle$ is large. Hence, a good measurement of this moment alone constrains $\bar{\Lambda}$ effectively.

Concerning the nonperturbative effects related to the $c\bar{c}$ loop, they have been calculated using HQET for the dilepton invariant mass spectrum away from the $c\bar{c}$ resonances [22]. We calculated the $1/m_c^2$ contributions to the hadronic moments using the amplitude given in [22]. The invariant mass and mixed moments give zero contribution for $m_s=0$. Thus, the corrections to the hadronic mass moments are vanishing, if we further neglect terms proportional to $(\lambda_2/m_c^2)\bar{\Lambda}$ and $(\lambda_2/m_c^2)\lambda_i$, with $i=1,2$. For the hadron energy moments we obtain numerically

$$\begin{aligned}\Delta\langle E_H\rangle_{1/m_c^2} &= m_B\Delta\langle x_0\rangle = -0.007\ \mathrm{GeV},\\ \Delta\langle E_H^2\rangle_{1/m_c^2} &= m_B^2\Delta\langle x_0^2\rangle = -0.013\ \mathrm{GeV}^2,\end{aligned} \tag{13}$$

leading to a correction of order -0.3% to the short-distance values presented in Table I.

Of course, the utility of the hadronic moments calculated above is only in conjunction with the experimental cuts. The optimal experiments cuts in $B\to X_s l^+ l^-$ remain to be defined, but for the cuts used by the CLEO Collaboration we have studied the effects in the HQET-like Fermi motion (FM) model [14]. We find that the hadronic moments in the HQET and FM model are very similar and CLEO-type cuts remove the bulk of the $c\bar{c}$ resonant contributions [11]. We hope to return to this and the related issue of doing an improved theoretical error estimate in the HQET context in a future publication. The power corrections presented here in the hadron spectrum and hadronic spectral moments in $B\to X_s l^+ l^-$ are the first results in this decay.

In summary, we have presented the results on the spectral hadronic moment $\langle E_H^n\rangle$ and $\langle S_H^n\rangle$ for $n=1,2$ and have worked out their dependence on the HQET parameters $\bar{\Lambda}$ and λ_1. The correlations in $B\to X_s l^+ l^-$ are shown to be different than the ones in the semileptonic decay $B\to X l \nu_l$. This complementarity allows, in principle, a powerful method to determine them precisely from data on $B\to X l \nu_l$ and $B\to X_s l^+ l^-$ in forthcoming high luminosity B facilities.

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