# Analysis of exclusive $B_s \to D_{s_0}(2317)\ell\bar{\nu}_\ell$ decay in "full" QCD



T. M. Aliev *†, M. Savcı ‡

Physics Department, Middle East Technical University, 06531 Ankara, Turkey

**Abstract**

The transition form factors of the semileptonic decay $B_s$ into scalar $D_{s_0}(2317)$ meson is calculated in the framework of 3–point QCD sum rule. The branching ratio is found to be $\sim 10^{-3}$ for the $B_s \to D_{s_0}(2317)\ell\bar{\nu}_\ell$ ($\ell = e, \mu$) decay, and $\sim 10^{-4}$ for the $B_s \to D_{s_0}(2317)\tau\bar{\nu}_\tau$ decay, respectively.

*e-mail: taliev@metu.edu.tr

†permanent address:Institute of Physics,Baku,Azerbaijan

‡e-mail: savci@metu.edu.tr

# 1 Introduction

Semileptonic decays of mesons containing charm and beauty quarks constitute a very important class of decays for determination of elements of the Cabibbo-Kobayashi-Maskawa (CKM) matrix, leptonic decay constants of heavy mesons, as well as for understanding the origin of CP violation which is related to the structure of the CKM matrix in the Standard Model (SM), because strong interactions involving semileptonic decays are more simple compared to that of hadronic decays. In semileptonic decays the long distance strong binding dynamics can be parametrized as transition form factors, calculation of which is the main problem of these decays. For estimation of transition form factors some nonperturbative approach is needed. Several methods, such as the quark model, lattice QCD, QCD sum rules, large energy and effective heavy quark theories, have been used to calculate the transition form factors. Among these approaches the QCD sum rules occupy a special place, since it is based on the fundamental QCD Lagrangian.

The QCD sum rules method [1] has been successfully applied to a wide variety of problems in hadron physics (see [2, 3] and references therein). The semileptonic decay $D \to \bar{K}^0 e \bar{\nu}_e$ is first studied in QCD sum rules with the three–point correlation function in [4]. Following this work, $D^+ \to K^0 e^+ \nu_e$, $D^+ \to K^{0*} e^+ \nu_e$ [5], $D \to \pi e \bar{\nu}_e$ [6], $D \to \rho e \bar{\nu}_e$ [7] and $B \to D(D^*) \ell \bar{\nu}_\ell$ [8], $D \to \phi \ell \bar{\nu}_\ell$ [9] are studied in the frame work of the same method. Note that $D_s \to \phi \ell \bar{\nu}_\ell$ decay is studied in light cone QCD sum rules [10], which is an alternative approach to the traditional QCD sum rules.

Transition form factors appearing in semileptonic decays depend not only on the dynamics of strong interactions between quarks in the initial and final state hadrons, but also on the structure of the hadrons involved in the semileptonic decays. Since in the present work we will consider a scalar meson in the final state, a few words about the scalar mesons are in order. The structure of the scalar mesons is still under debate. At present, there are different proposals about the nature of the scalar mesons that have been put forward, for example, their structures are considered as composed of $\bar{q}q$, multiquark $\bar{q}q\bar{q}q$ or meson–meson bound states.

For studying the structure of scalar mesons much more experimental data and theoretical analysis are needed. The observation of two narrow resonances with charm and strangeness $D_{s_0}(2317)$ in the $D_s\pi$ invariant mass distribution [11–17] and $D_{s_J}(2460)$ in the $D_s^*\pi^0$ and $D_s\gamma$ mass distributions [12–14, 17–19] have raised discussion about the nature and quark content on these states [20]. The natural identification consists in considering these states as the scalar $D_{s_0}(2317)$ and axial vector $D_{s_J}(2460)$ $\bar{c}s$ mesons, respectively. The result of the analysis of the radiative decays $D_{s_0}(2317) \to D_s^*\gamma$, $D_{s_J}^*(2460) \to D_s^*\gamma$ and $D_{s_J}^*(2460) \to D_{s_0}(2317)\gamma$ is in favor of the interpretation of quark content of these mesons as being ordinary $\bar{c}s$ mesons [21]. It is further observed that deviation of the transition amplitude in the infinite heavy quark limit compared to that of the finite quark mass case is quite sizeable. In the light of the $D_{s_0} \to D_s^*\gamma$ case, one can ask whether a change in the heavy quark effective theory results occurs for the $B_s \to D_{s_0}(2317)\ell\bar{\nu}_\ell$ decay for the finite quark mass, or not. In the present we try to answer this question. Note that this decay has been studied within the frame work of the heavy quark effective theory in [22].

In this work, we will study the semileptonic decay of $B_s$ meson to $D_{s_0}(2317)$ meson, i.e., $B_s \to D_{s_0}(2317)\ell\bar{\nu}_\ell$, in the frame work of QCD sum rules method. The paper is organized as follows: In section 2, we derive the sum rules for the transition form factors. Section 3 is devoted to the numerical analysis and discussion, and contains a summary of our results and conclusions.

# 2 Sum rules for the $B_s \to D_{s_0}(2317)$ transition form factors

The amplitude of the $B_s \to D_{s_0}(2317)\ell\bar{\nu}_\ell$ decay can be presented in the following form

$$\mathcal{M} = \frac{G_F}{\sqrt{2}} V_{cb} \bar{\nu}_\ell \gamma_\mu (1-\gamma_5)\ell \left\langle D_{s_0} \left| \bar{c}\gamma^\mu (1-\gamma_5) b \right| B_s \right\rangle \ , \tag{1}$$

where $V_{cb}$ is the CKM matrix element which describes the transition of a $b$ quark into a $c$ quark, and for the sake of simplicity, in all further calculations we will denote $D_{s_0}(2317)$ as $D_0$. The main problem related to the calculation of the matrix element $\langle D_0 | \bar{c}\gamma^\mu(1-\gamma_5)b | B_s\rangle$. Obviously, the vector part of weak $\bar{c}\gamma^\mu(1-\gamma_5)b$ current does not contribute to the matrix element considered above, which immediately follows from parity property of the hadrons and weak current; and only axial part of the weak current gives nonzero contribution. From the Lorentz invariance, this matrix element can be parametrized in terms of the form factors in the following way:

$$\langle D_0(p') | \bar{c}\gamma_\mu\gamma_5 b | B_s(p)\rangle = i\Big[f_+ \mathcal{P}_\mu + f_- q_\mu\Big] \ , \tag{2}$$

where $f_+(q^2)$ and $f_-(q^2)$ are the transition form factors, $\mathcal{P}_\mu = (p+p')_\mu$ and $q_\mu = (p-p')_\mu$. It is well known that $f_-$ is proportional to the lepton mass, and especially for the $\tau$ case it can substantially be important. For this reason we will take both of the form factors into consideration. In calculation of the form factors $f_+(q^2)$ and $f_-(q^2)$ we will employ the QCD sum rules method and proceed by considering the following correlator:

$$\Pi_\mu(p^2, p'^2, q^2) = i^2 \int d^4x d^4y e^{i(p'y - px)} \left\langle 0 \left| J_{D_0}(y) J_\mu^A(0) J_5(x) \right| 0 \right\rangle \ , \tag{3}$$

where $J_{D_0}(y) = \bar{s}c$, $J_5 = \bar{s}\gamma_5 b$ and $J_\mu^A = \bar{c}\gamma_\mu\gamma_5 b$ are the interpolating currents of the scalar $D_0$, $B_s$ mesons and weak axial currents, respectively.

Let us first calculate the phenomenological part of the correlator given in Eq. (3). This can be obtained by inserting the complete set of intermediate states with the same quantum number as the currents $J_{D_0}$ and $J_5$. isolating the pole terms of the lowest scalar and pseudoscalar $D_0$ and $B_s$ mesons, we get the following representation of the above–mentioned correlator

$$\begin{aligned}\Pi_\mu(p^2, p'^2, q^2) &= \frac{\langle 0 | J_{D_0} | D_0\rangle \langle D_0 | \bar{c}\gamma_\mu\gamma_5 b | B_s\rangle \langle B_s | \bar{s}\gamma_5 b | 0\rangle}{(m_{D_0}^2 - p'^2)(m_{B_s}^2 - p^2)} \\ &+ \sum_h \frac{\langle 0 | J_{D_0} | h(p')\rangle \langle h(p') | \bar{c}\gamma_\mu\gamma_5 b | \mathcal{H}(p)\rangle \langle \mathcal{H}(p) | \bar{s}\gamma_5 b | 0\rangle}{(p'^2 - m_h^2)(p^2 - m_{\mathcal{H}}^2)} \ .\end{aligned} \tag{4}$$

The second term in Eq. (4) takes into account higher states and continuum contributions; and $h$ and $\mathcal{H}$ form a complete set of mesons having the same quantum numbers as the ground state mesons.

The matrix elements in Eq. (4) are defined in the standard way as:

$$\begin{aligned}
\langle 0 \,| J_{D_0} |\, D_0 \rangle &= f_{D_0} m_{D_0} \ , \\
\langle B_s \,| \bar{s}\gamma_5 b |\, 0 \rangle &= -i \frac{f_{B_s} m_{B_s}^2}{m_b + m_s} \ ,
\end{aligned} \tag{5}$$

where $f_{D_0}$ and $f_{B_s}$ are the leptonic decay constants of $D_0$ and $B_s$ mesons, respectively. Using (5), Eq. (4) can be written as

$$\Pi_\mu(p^2, p'^2, q^2) = -\frac{f_{B_s} m_{B_s}^2}{(m_b + m_s)} \frac{f_{D_0} m_{D_0}}{(m_{D_0}^2 - p'^2)(m_{B_s}^2 - p^2)} \Big[ f_+ \mathcal{P}_\mu + f_- q_\mu \Big] + \text{excited states} \ . \tag{6}$$

In accordance with the QCD sum rules philosophy, $\Pi_\mu(p^2, p'^2, q^2)$ can be calculated from QCD side wit the help of the operator product expansion method (OPE) in the deep Euclidean region $p^2 \ll (m_b + m_c)^2$ and $p'^2 \ll (m_c + m_s)^2$. Equating the two different representations of $\Pi_\mu$ gives us sum rules for the form factors $f_+(q^2)$ and $f_-(q^2)$.

The theoretical part of the correlator is calculated by means of OPE, and up to operators having dimension d=6, it is determined by the bare–loop and the power corrections (Figs. (1)–(3)), from the operators with d=3 $\langle \bar\psi \psi \rangle$, d=5 $m_0^2 \langle \bar\psi \psi \rangle$ and d=6 $\langle \bar\psi \psi \rangle^2$. Our calculation shows that d=4 operator $\left\langle G_{\mu\nu}^2 \right\rangle$ gives very small contribution, and for this reason we do not consider it in the present work.

In calculating the bare–loop contribution, we we first write the double dispersion representation as

$$f_i^{per} = -\frac{1}{(2\pi)^2} \int ds ds' \frac{\rho_i(s, s', q^2)}{(s - p^2)(s' - p'^2)} + \text{subtraction terms} \ . \tag{7}$$

The spectral density $\rho_i(s, s', q^2)$ can be calculated from the usual Feynman integral with the help of Gutkovsky rule, i.e., by replacing the denominators of the propagators as follows:

$$\frac{1}{p^2 - m^2} \to -2\pi i \delta(p^2 - m^2) \ ,$$

which implies that all quarks are real.

After standard calculations for the spectral densities we obtain:

$$\begin{aligned}
\rho_+(s, s', q^2) &= \frac{1}{4\lambda^{1/2}(s, s', q^2)} \Big\{ (\Delta' + \Delta)(1 + A + B) + [(m_b + m_c)^2 - q^2](A + B) \\
&+ 2m_s(m_b - m_c)(1 + A + B) \Big\} \\
\rho_-(s, s', q^2) &= \frac{1}{4\lambda^{1/2}(s, s', q^2)} \Big\{ [\Delta' + \Delta + (m_b + m_c)^2 - q^2 + 2m_s(m_b - m_c)](A - B) \\
&+ \Delta' - \Delta - 2m_s(m_b + m_c) \Big\} \ ,
\end{aligned} \tag{8}$$

where $\Delta' = s - m_c^2$ and $\Delta = s - m_b^2$, and

$$\begin{aligned}
A &= \frac{1}{\lambda(s, s', q^2)} [-(s + s' - q^2)\Delta' + 2\Delta s'] \ , \\
B &= \frac{1}{\lambda(s, s', q^2)} [-(s + s' - q^2)\Delta + 2\Delta' s] \ ,
\end{aligned}$$

and $\lambda(s,s',q^2) = s^2 + s'^2 + q^4 - 2sq^2 - 2s'q^2 - 2ss'$. Here and in all following expressions, subscripts $+$ and $-$ correspond to the coefficients of the structures proportional to $\mathcal{P}_\mu$ and $q_\mu$, respectively. Note that, in deriving Eqs. (7) and (8) we retain only linear terms in $m_s$ in order to take SU(3) violation effects into account, and higher order $m_s$ terms are all neglected. The integration region in for the perturbative contribution in Eq. (7) is determined from the condition that arguments of the three $\delta$ functions might vanish simultaneously. The physical region in $s$ and $s'$ plane is described by the following inequalities:

$$-1 \le \frac{2ss' - 2m_c^2 s + (s+s'-q^2)(m_b^2 - s)}{\lambda^{1/2}(s,s',q^2)(m_b^2-s)} \le +1 \ .$$

According to the quark–hadron duality, the contribution of higher states in phenomenological part is parametrized in correspondence with the spectral density starting from $s > s_0$ and $s' > s_0'$.

In what follows, we present the contributions of d=3, 5 and 6 operators.

$$f_+^{(3)} = \frac{1}{2}\langle \bar{s}s\rangle \frac{m_b - m_c}{rr'} \ , \tag{9}$$

$$f_-^{(3)} = -\frac{1}{2}\langle \bar{s}s\rangle \frac{m_b + m_c}{rr'} \ , \tag{10}$$

$$f_+^{(4)} = \frac{1}{4} m_s \langle \bar{s}s\rangle \left[ -\frac{m_c(m_b-m_c)}{rr'^2} + \frac{m_b(m_b-m_c)}{r^2 r'} - \frac{2}{rr'} \right] , \tag{11}$$

$$f_-^{(4)} = \frac{1}{4} m_s (m_b+m_c) \langle \bar{s}s\rangle \left[ \frac{m_c}{rr'^2} - \frac{m_b}{r^2 r'} \right] , \tag{12}$$

$$\begin{aligned} f_+^{(5)} = & -\frac{1}{12} m_0^2 \langle \bar{s}s\rangle \left[ \frac{3m_c^2(m_b-m_c)}{rr'^3} + \frac{3m_b^2(m_b-m_c)}{r^3 r'} + \frac{2(m_b-2m_c)}{rr'^2} \right. \\ & \left. + \frac{2(2m_b-m_c)}{r^2 r'} + \frac{(m_b-m_c)(2m_b^2+m_b m_c+2m_c^2-2q^2)}{r^2 r'^2} \right] , \end{aligned} \tag{13}$$

$$\begin{aligned} f_-^{(5)} = & \frac{1}{12} m_0^2 \langle \bar{s}s\rangle \left[ \frac{3m_c^2(m_b+m_c)}{rr'^3} + \frac{3m_b^2(m_b+m_c)}{r^3 r'} + \frac{2(m_b+3m_c)}{rr'^2} \right. \\ & \left. + \frac{2(m_c+3m_b)}{r^2 r'} + \frac{(m_b+m_c)(2m_b^2+m_b m_c+2m_c^2-2q^2)}{r^2 r'^2} \right] , \end{aligned} \tag{14}$$

$$\begin{aligned} f_+^{(6)} = & \frac{4}{81}\pi\alpha_s \langle \bar{s}s\rangle^2 \left[ \frac{12m_c^3(m_b-m_c)}{rr'^4} - \frac{12m_b^3(m_b-m_c)}{r^4 r'} \right. \\ & + \frac{4m_c(m_b-m_c)(2m_b^2+m_b m_c+2m_c^2-2q^2)}{r^2 r'^3} + \frac{8m_c(7m_b-8m_c)}{rr'^3} \\ & - \frac{4m_b(m_b-m_c)(2m_b^2+m_b m_c+2m_c^2-2q^2)}{r^3 r'^2} - \frac{8m_b(8m_b-7m_c)}{r^3 r'} \\ & \left. - \frac{4(5m_b^2-20m_b m_c+5m_c^2+2q^2)}{r^2 r'^2} + \frac{48}{rr'^2} + \frac{48}{r^2 r'} \right] \\ & + \frac{1}{9} m_0^2 m_s \langle \bar{s}s\rangle \left[ -\frac{6m_c^3(m_b-m_c)}{rr'^4} - \frac{m_c(m_b-m_c)(5m_b^2+4m_b m_c+5m_c^2-5q^2)}{r^2 r'^3} \right. \\ & - \frac{2m_c(7m_b-11m_c)}{rr'^3} - \frac{m_b(m_b-m_c)(m_b^2+8m_b m_c+m_c^2-q^2)}{r^3 r'^2} \end{aligned}$$

$$+\frac{5m_b^2-20m_bm_c+11m_c^2-4q^2}{r^2r'^2}+\frac{6m_b^3(m_b-m_c)}{r^4r'}+\frac{2m_b(5m_b-13m_c)}{r^3r'}\Bigg], \quad (15)$$

$$f_-^{(6)} = \frac{4}{81}\pi\alpha_s\langle\bar{s}s\rangle^2\Bigg[\frac{-12m_c^3(m_b+m_c)}{rr'^4}+\frac{12m_b^3(m_b+m_c)}{r^4r'}$$
$$-\frac{4m_c(m_b+m_c)(2m_b^2+m_bm_c+2m_c^2-2q^2)}{r^2r'^3}-\frac{8m_c(7m_b+9m_c)}{rr'^3}$$
$$+\frac{4m_b(m_b+m_c)(2m_b^2+m_bm_c+2m_c^2-2q^2)}{r^3r'^2}+\frac{8m_b(9m_b+7m_c)}{r^3r'}$$
$$+\frac{28(m_b^2-m_c^2)}{r^2r'^2}+\frac{8}{rr'^2}-\frac{8}{r^2r'}\Bigg]$$
$$+\frac{1}{9}m_0^2m_s\langle\bar{s}s\rangle\Bigg[\frac{6m_c^3(m_b+m_c)}{rr'^4}+\frac{m_c(m_b+m_c)(5m_b^2+4m_bm_c+5m_c^2-5q^2)}{r^2r'^3}$$
$$+\frac{2m_c(7m_b+12m_c)}{rr'^3}+\frac{m_b(m_b+m_c)(m_b^2+8m_bm_c+m_c^2-q^2)}{r^3r'^2}$$
$$-\frac{m_b^2-18m_bm_c-7m_c^2}{r^2r'^2}-\frac{6m_b^3(m_b+m_c)}{r^4r'}+\frac{4}{rr'^2}-\frac{4}{r^2r'}-\frac{2m_b(12m_b+13m_c)}{r^3r'}\Bigg], \quad (16)$$

where $r=p^2-m_b^2$, $r'=p'^2-m_c^2$. Note that, in the present work we neglect the $\mathcal{O}(\alpha_s)$ corrections to the bare loop. For consistency, we also neglect $\mathcal{O}(\alpha_s)$ corrections in determination of the leptonic decay constants $f_{B_s}$ and $f_{D_0}$.

Substitute Eqs. (8)–(16) into Eq. (7) and applying double Borel transformation $\hat{\mathcal{B}}$ with respect to the variables $p^2$ and $p'^2$ ($p^2\to M_1^2,\ p'^2\to M_2^2$) in order to suppress the contributions of higher states and continuum, we get the following sum rules for the form factors $f_+$ and $f_-$:

$$f_\pm(q^2) = -\frac{(m_b+m_c)}{f_{B_s}m_{B_s}^2}\frac{1}{f_{D_s}m_{D_S}}e^{(m_{B_s}^2/M_1^2+m_{D_s}^2/M_2^2)}\Bigg[\hat{B}\Big(f_\pm^{(3)}+f_\pm^{(4)}+f_\pm^{(5)}+f_\pm^{(6)}\Big)$$
$$-\frac{1}{(2\pi)^2}\int dsds'\rho_\pm(s,s',q^2)e^{-s/M_1^2-s'/M_2^2}\Bigg]. \quad (17)$$

Double Borel transformation is implemented by the following expression:

$$\hat{\mathcal{B}}\frac{1}{r^m}\frac{1}{r'^n}\to(-1)^{m+n}\frac{1}{\Gamma(m)}\frac{1}{\Gamma(n)}e^{-s/m_b^2}e^{-s'/m_c^2}\ .$$

# 3 Numerical analysis

In this section we present our numerical analysis for the form factors $f_+(q^2)$ and $f_-(q^2)$. It follows from the expressions of these form factors that the main input parameters needed are the condensates, leptonic decay constants of $B_s$ and $D_0$ mesons, continuum thresholds $s_0$ and $s_0'$ and Borel parameters $M_1^2$ and $M_2^2$.

In further numerical analysis we choose the value of the condensates at a fixed renormalization scale of about 1 $GeV$. The values of the condensates are taken from [22] and can be listed as follows:

$$\langle \bar\psi\psi\rangle\Big|_{\mu=1~GeV} = -(240\pm 10~MeV)^3~,$$
$$\langle \bar s s\rangle = (0.8\pm 0.1)\langle\bar\psi\psi\rangle~.$$

The quark masses are taken to be $m_c = (\mu = m_c) = 1.275 \pm 0.015~GeV$, $m_s(1~GeV) \simeq 142~MeV$ [22], $m_b = (4.7 \pm 0.1)~GeV$ [23]. For the values of the leptonic decay constants of $B_s$ and $D_0$ mesons we use the results obtained from two–point QCD analysis: $f_{B_s} = 209 \pm 38~MeV$ [3] and $f_{D_0} = 225 \pm 25~MeV$ [21]. The threshold parameters $s_0$ and $s_0'$ are also determined from the two point QCD sum rules: $s_0 = (35 \pm 2)~GeV^2$ [3] and $s_0' = (2.5~GeV)^2$ [21]. The Borel parameters $M_1^2$ and $M_2^2$ are auxiliary quantities and therefore the results of physical quantities should not depend on them, if OPE can be calculated up to all order. In QCD sum rule method, OPE is truncated at some finite order. For this reason, we need to choose "working" regions for the Borel parameters where form factors are supposed to be practically independent of them. The choice of the working region for the Borel parameters $M_1^2$ and $M_2^2$ should be based, on the one side, on the condition that that the continuum contribution should be small, and on the other side, the convergence of the power corrections. As a result of the above–mentioned conditions, the best stability is achieved for $10~GeV^2 \le M_1^2 \le 15~GeV^2$ and $4~GeV^2 \le M_2^2 \le 7~GeV^2$.

As a result of the above–summarized considerations, our analysis leads to th following predictions for the form factors at $q^2 = 0$:

$$\begin{aligned} f_+ &= 0.20 \pm 0.05~, \\ f_- &= -0.32 \pm 0.08~. \end{aligned} \tag{18}$$

The errors are can be attributed to the variation of the thresholds, decay constants, uncertainties in condensates and in quark masses.

For completeness, we also present the results of HQET for the above–mentioned form factors that predicts, $f_+ = -0.37$ and $f_- = -0.15$. Indeed, we observe that finite quark mass effects are essential in determination of the form factors.

In order to estimate the width of $B_s \to D_0 \ell\bar\nu_\ell$ it is necessary to know the $q^2$ dependence of the form factors $f_+(q^2)$ and $f_-(q^2)$ in the whole physical region $m_\ell^2 \le q^2 \le (m_{B_s} - m_{D_0})^2$. The $q^2$ dependence of the form factors can be calculated from QCD sum rules (for details, see [5, 6]). In the present work we have analyzed this dependence and used it in our numerical calculations.

Having decided on the parametrization of the form factors, it is not difficult to obtain the expression for the differential decay rate

$$\begin{aligned} \frac{d\Gamma}{dq^2} &= \frac{1}{192\pi^3 m_{B_s}^3} G^2 \left|V_{cb}\right|^2 \lambda^{1/2}(m_{B_s}^2, m_{D_0}^2, q^2) \left(\frac{q^2 - m_\ell^2}{q^2}\right)^2 \\ &\times \Bigg\{ -\frac{(2q^2 + m_\ell^2)}{2} \Big[ \left|f_+(q^2)\right|^2 (2m_{B_s}^2 + 2m_{D_0}^2 - q^2) + 2(m_{B_s}^2 - m_{D_0}^2)\mathrm{Re}[f_+(q^2) f_-^*(q^2)] \\ &+ \left|f_-(q^2)\right|^2 q^2 \Big] + \frac{(q^2 + 2m_\ell^2)}{q^2} \Big[ \left|f_+(q^2)\right|^2 (m_{B_s}^2 - m_{D_0}^2)^2 \\ &+ 2(m_{B_s}^2 - m_{D_0}^2) q^2 \mathrm{Re}[f_+(q^2) f_-^*(q^2)] + \left|f_-(q^2)\right|^2 q^4 \Big] \Bigg\}~. \end{aligned}$$

Taking into account the $q^2$ dependence of the form factors $f_+$ and $f_-$, and performing integration over $q^2$ and using the total life–time $\tau_{B_s} = 1.46 \times 10^{-12}\ s$ [24], we get for the branching ratio

$$\begin{aligned}
\mathcal{B}(B_s \to D_0 \ell \bar{\nu}_\ell) &\simeq 10^{-3}\ , \quad (\ell = e,\ \mu)\ , \\
\mathcal{B}(B_s \to D_0 \tau \bar{\nu}_\tau) &\simeq 10^{-4}\ .
\end{aligned}$$

Our predictions on the branching ratio is considerably smaller compared to that of the HQET results.

At the end of our analysis, we would like to note that, although current $B$ meson factories do not produce $B_s$ meson, it is hoped to be possible to study the weak decays of $B_s$ mesons in future–planned experiments at LHC. The exclusive $B_s \to D_0 \ell \bar{\nu}_\ell$ decay can be studied via the resonant production of the scalar $D_0$ meson in the weak decay of the $B_s$ meson. This analysis can give valuable essential information about the quark content of the scalar $D_0$ meson.

In summary, we study the semileptonic $B_s \to D_0 \ell \bar{\nu}_\ell$ decay in the framework of 3–point QCD sum rules. We calculate the transition form factors, and using these predictions, we estimate the branching ratio of the $B_s \to D_0 \ell \bar{\nu}_\ell$ decay.

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