Semileptonic $B_s \to D_{s2}^*(2573)\ell\bar{\nu}_\ell$ transition in QCD

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Abstract

We analyze the semileptonic $B_s \to D_{s2}^*(2573)\ell\bar{\nu}_\ell$ transition, where $\ell = \tau$, $\mu$ or $e$, within the standard model. We apply the QCD sum rule approach to calculate the transition form factors entering the low energy Hamiltonian defining this channel. The fit functions of the form factors are used to estimate the total decay widths and branching fractions in all lepton channels. The orders of branching ratios indicate that this transition is accessible at LHCb in near future.

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1 Introduction

The semileptonic $B$ meson decay channels are known as useful tools to accurately calculate the Standard Model (SM) parameters like determination of the Cabibbo-Kobayashi-Maskawa (CKM) quark mixing matrix, check the validity of the SM, describe the origin of the CP violation and search for new physics effects. By recent experimental progresses, it has become precise measurements available, and it is possible to perform precision calculations. Although the $B$ meson decays are studied efficiently both theoretically and experimentally (see for instance [1–11]), most of $B_s$ properties are not very clear yet (for some related theoretical and experimental studies on this meson see [12–21] and references therein). Since the detection and identification of this heavy meson is relatively difficult in the experiment, the theoretical and phenomenological studies on the spectroscopy and decay properties of this mesons can play essential role in our understanding of its non-perturbative dynamics, calculating the related parameters of the SM and providing opportunities to search for possible new physics contributions.

In the literature, there are a lot of theoretical studies devoted to the semileptonic transition of $B_s$ into the pseudoscalar $D_s$ and vector $D_s^*$ charmed-strange mesons. But, we have no study on the semileptonic transitions of this meson into the tensor charmed-strange meson in final state, although it is expected to have considerable contribution to the total decay width of the $B_s$ meson. In this accordance, in the present study, we investigate the semileptonic $B_s \rightarrow D_{s2}^*(2573)\ell\bar{\nu}_\ell$ transition in the framework of three-point QCD sum rule [22] as one of the most attractive and powerful techniques in hadron phenomenology, where the $D_{s2}^*(2573)$ is the low lying charmed-strange tensor meson with $J^P = 2^+$. In particular, we calculate the transition form factors entering the low energy matrix elements defining the transition under consideration. We find the working regions of the auxiliary parameters entering to calculations from different transformations, considering the criteria of the method used. This is followed by finding the behavior of the form factors in terms of the transferred momentum squared, which are then used to estimate the total width and branching fraction in all lepton channels. Note that the semileptonic $B \rightarrow D_s^*(2460)\ell\bar{\nu}_\ell$ decay channel is analyzed in [23] using the same method. The spectroscopic properties of the charmed-strange tensor meson $D_{s2}^*(2573)$ is also investigated in [24] using a two-point correlation function.

The layout of the paper is as follows. In next section, the QCD sum rules for the four form factors relevant to semileptonic $B_s \rightarrow D_{s2}^*(2573)\ell\bar{\nu}_\ell$ transition are obtained. Section 3 contains numerical analysis of the form factors, calculation of their behavior in terms of $q^2$ as well as the estimation of the total decay width and branching ratio for the transition under consideration.

2 Theoretical framework

In order to calculate the form factors, associated with the semileptonic $B_s \rightarrow D_{s2}^*(2573)\ell\bar{\nu}_\ell$ transition via QCD sum rule formalism, we consider the following three-point correlation function:
\[ \Pi_{\mu\alpha\beta} = i^2 \int d^4x \int d^4y e^{-ip\cdot x} e^{ip'\cdot y} \langle 0 \mid \mathcal{T} \left[ J_{\alpha\beta}^{D_2^*(2573)}(y) J_{\mu}^{tr}(0) J_{B_s}^{\dagger}(x) \right] \mid 0 \rangle, \]  

where \( \mathcal{T} \) is the time ordering operator and \( J_{\mu}^{tr}(0) = \bar{c}(0)\gamma_\mu(1 - \gamma_5)b(0) \) is the transition current. The interpolating currents of the \( B_s \) and \( D_2^*(2573) \) mesons can be written in terms of the quark fields as

\[ J_{B_s} = \bar{s}(x)\gamma_5b(x), \]  

and

\[ J_{D_2^*(2573)}^{\alpha\beta}(y) = \frac{i}{2} \left[ \bar{s}(y)\gamma_\alpha \overset{\leftrightarrow}{D}_\beta(y)c(y) + \bar{s}(y)\gamma_\beta \overset{\leftrightarrow}{D}_\alpha(y)c(y) \right]. \]  

Here \( \overset{\leftrightarrow}{D}_\beta(y) \) is the covariant derivative that acts on the left and right, simultaneously. It is given as

\[ \overset{\leftrightarrow}{D}_\beta(y) = \frac{1}{2} \left[ \overset{\rightarrow}{D}_\beta(y) - \overset{\leftarrow}{D}_\beta(y) \right], \]  

with

\[ \overset{\rightarrow}{D}_\beta(y) = \overset{\to}{D}_\beta(y) - ig2\lambda^a A^a_\beta(y), \]  

\[ \overset{\leftarrow}{D}_\beta(y) = \overset{\leftarrow}{\partial}_\beta(y) + ig2\lambda^a A^a_\beta(y), \]  

where \( \lambda^a \) and \( A^a_\beta(x) \) denote the Gell-Mann matrices and the external gluon fields, respectively.

According to the method used, in order to find the QCD sum rules for transition form factors, we shall calculate the aforesaid correlation function, once in terms of hadronic parameters and the second in terms of QCD parameters making use of operator product expansion (OPE). By equating these two representations to each other through a dispersion relation, we obtain the sum rules for form factors. To stamp down the contributions of the higher states and continuum, a double Borel transformation with respect to the \( p^2 \) and \( p'^2 \) is performed on both sides of the sum rules obtained and the quark-hadron duality assumption is used.

2.1 The hadronic representation

In order to calculate the hadronic side of the correlator in Eq.(1), we insert two complete sets of the initial \( B_s \) and the final \( D_2^*(2573) \) states with the same quantum numbers as the interpolating currents into the correlator. After performing four-integrals over \( x \) and \( y \), we obtain

\[ \Pi^{\text{had}}_{\mu\alpha\beta} = \frac{\langle 0 \mid J_{\alpha\beta}^{D_2^*(2573)} \mid D_2^*(2573)(p', \epsilon) \rangle \langle D_2^*(2573)(p', \epsilon) \mid J_{\mu}^{tr}(0) \mid B_s(p) \rangle \langle B_s(p) \mid J_{B_s}^{\dagger} \mid 0 \rangle}{(p^2 - m^2_{B_s})(p'^2 - m^2_{D_2^*(2573)})} \]  

\[ + \ldots, \]  

(6)
where $\cdots$ represents contributions of the higher states and continuum, and $\epsilon$ is the polarization tensor of the $D_{s2}^*(2573)$ tensor meson. We can parameterize the matrix elements appearing in the above equation in terms of decay constants, masses and form factors as

\[
\langle 0 \mid J_{a\beta}^{D_{s2}^*(2573)} \mid D_{s2}^{*}(2573)(p', \epsilon) \rangle = m_{D_{s2}^*(2573)}^3 f_{D_{s2}^*(2573)}(2573)\epsilon_{a\beta},
\]

\[
\langle B_s(p) \mid J_{\mu}^{\gamma} \mid 0 \rangle = -i \frac{f_{B_s} m_{B_s}^2}{m_s + m_b},
\]

\[
\langle D_{s2}^{*}(2573)(p', \epsilon) \mid J_{\mu}^{tr}(0) \mid B_s(p) \rangle = h(q^2)\epsilon_{\mu\nu\alpha\beta} \epsilon^{\nu\lambda} P^{\lambda} P_{\alpha} q_{\beta} - iK(q^2)\epsilon^{\nu\mu} P^{\nu}
\]

\[
- i\epsilon_{\alpha\beta} P^{\alpha} P^{\beta} [P_{\mu} b_{+}(q^2) + q_{\mu} b_{-}(q^2)],
\]

where $q = p - p'$, $P = p + p'$; and $h(q^2)$, $K(q^2)$, $b_{+}(q^2)$ and $b_{-}(q^2)$ are transition form factors.

Now, we combine Eqs. (6) and (7) and performing summation over the polarization tensors via

\[
\epsilon_{\alpha\beta}\epsilon^{\nu\theta} = \frac{1}{2} T_{\alpha\nu} T_{\beta\theta} + \frac{1}{2} T_{\alpha\theta} T_{\beta\nu} - \frac{1}{3} T_{\alpha\beta} T_{\nu\theta},
\]

where

\[
T_{\alpha\nu} = -g_{\alpha\nu} + \frac{p_{\alpha}' p_{\nu}'}{m_{D_{s2}^*(2573)}}.
\]

This procedure brings us to the final representation of the hadronic side, viz.

\[
\Pi_{\mu\alpha\beta}^{had} = \frac{f_{D_{s2}^{*}} f_{B_s} m_{D_{s2}^*(2573)} m_{B_s}^2}{8(m_b + m_s)(p^2 - m_{B_s}^2)(p'^2 - m_{D_{s2}^*(2573)}^2)} \left\{ \frac{2}{3} \left[ -\Delta K(q^2) + \Delta' b_{+}(q^2) \right] q_{\mu} g_{\beta\alpha} + \frac{2}{3} \left[ (\Delta - 4m_{D_{s2}^*(2573)}^2)K(q^2) + \Delta' b_{+}(q^2) \right] P_{\mu} g_{\beta\alpha} + i(\Delta - 4m_{D_{s2}^*(2573)}^2)h(q^2)\epsilon_{\nu\beta\alpha} P_{\lambda} q_{\nu}
\right.
\]

\[
+ \Delta K(q^2) q_{\alpha} g_{\beta\mu} + \text{other structures} \right\} + \ldots,
\]

where

\[
\Delta = m_{B_s}^2 + 3m_{D_{s2}^*(2573)}^2 - q^2,
\]

and

\[
\Delta' = m_{B_s}^4 - 2m_{B_s}^2 (m_{D_{s2}^*(2573)}^2 + q^2) + (m_{D_{s2}^*(2573)}^2 - q^2)^2.
\]

### 2.2 The OPE representation

The OPE side of the correlation function is calculated in deep Euclidean region. For this aim, we insert the explicit forms of the interpolating currents into the correlation function in Eq. (1). After performing contractions via the Wick’s theorem, we obtain the following result in terms of the heavy and light quarks propagators:

\[
\Pi_{\mu\alpha\beta}^{OPE} = \frac{-i^3}{2} \int d^4x \int d^4y e^{-ip \cdot x} e^{ip' \cdot y} 
\]

\[
\times \left\{ Tr \left[ S_{s}^{ki}(x - y) \gamma_{\alpha} \bar{D}_{\beta}(y) S_{c}^{ij}(y) \gamma_{\mu}(1 - \gamma_5) S_{b}(-x) j_{k} \gamma_5 \right] + [\beta \leftrightarrow \alpha] \right\}.
\]

(13)
The heavy and light quarks propagators appearing in above equation and up to terms taken into account in the calculations are given by

\begin{equation}
S_Q^{ij}(x) = \frac{i}{(2\pi)^4} \int d^4k e^{-ik\cdot x} \left\{ \frac{\delta_{ij}}{k - m_Q} - \frac{g_s G_{ij}^a}{4} \frac{\sigma_{a\beta}(k + m_Q) + (k + m_Q)\sigma_{a\beta}}{(k^2 - m_Q^2)^2} \right. \\
+ \left. \delta_{ij} \frac{\alpha_s}{3} \frac{G^{ij}_G m_Q^2 + m_Q^2 k^2}{(k^2 - m_Q^2)^4} + \cdots \right\},
\end{equation}

where \( Q = b \) or \( c \), and

\begin{equation}
S_s^{ij}(x) = \frac{i}{2\pi^2 x^4} \delta_{ij} - \frac{m_s}{4\pi^2 x^2} \delta_{ij} - \frac{\langle ss \rangle}{12} \left( 1 - i \frac{m_s}{4} \frac{\not{x}}{x} \right) \delta_{ij} - \frac{x^2}{192} \frac{m_s^2 \langle ss \rangle}{x} \left( 1 - i \frac{m_s}{6} \frac{\not{x}}{x} \right) \delta_{ij} \\
- \frac{ig_s G^{ij}_G}{32\pi^2 x^2} \left[ \not{x} \sigma^{\eta} + \sigma^{\eta} \not{x} \right] + \cdots.
\end{equation}

To proceed, we insert the expressions of the heavy and light propagators into Eq. (13) and perform the derivatives with respect to \( x \) and \( y \). Then, we transform the calculations to the momentum space and make the \( x_\mu \to i\frac{\partial}{\partial y_\mu} \) and \( y_\mu \to -i\frac{\partial}{\partial y_\mu} \) replacements. We perform the two four-integrals coming from the heavy quark propagators with the help of two Dirac delta functions appearing in the calculations. Finally, we perform the last four-integral using the Feynman parametrization, viz.

\begin{equation}
\int dt_i \frac{(t_i^2)^{\beta}}{(t_i^2 + L)^{\alpha}} = \frac{i\pi^2(-1)^{\beta-\alpha} \Gamma(\beta + 2) \Gamma(\alpha - \beta - 2)}{\Gamma(2) \Gamma(\alpha) [-L]^{\alpha-\beta-2}}.
\end{equation}

Eventually, we get the OPE side of the three-point correlation function in terms of the selected structures and the perturbative and non-perturbative parts as

\begin{equation}
\Pi_{\mu\alpha}^{\text{OPE}} = \left( \Pi_1^{\text{pert}}(q^2) + \Pi_1^{\text{non-pert}}(q^2) \right) q_\alpha g_{\beta\mu} + \left( \Pi_2^{\text{pert}}(q^2) + \Pi_2^{\text{non-pert}}(q^2) \right) q_\mu g_{\beta\alpha} \\
+ \left( \Pi_3^{\text{pert}}(q^2) + \Pi_3^{\text{non-pert}}(q^2) \right) P_\mu g_{\beta\alpha} + \left( \Pi_4^{\text{pert}}(q^2) + \Pi_4^{\text{non-pert}}(q^2) \right) \varepsilon_{\lambda\nu\beta\mu} P_\lambda P_\alpha q_\nu \\
+ \text{other structures},
\end{equation}

where the perturbative parts \( \Pi_i^{\text{pert}}(q^2) \) can be written in terms of the double dispersion integrals as

\begin{equation}
\Pi_i^{\text{pert}}(q^2) = \int ds \int ds' \frac{\rho_i(s, s', q^2)}{(s - p^2)(s' - p'^2)}. 
\end{equation}

The spectral densities \( \rho_i(s, s', q^2) \) are given by the imaginary parts of the \( \Pi_i^{\text{pert}}(q^2) \) functions, i.e., \( \rho_i(s, s', q^2) = \frac{1}{\pi} Im[\Pi_i^{\text{pert}}(q^2)] \). After lengthy calculations, the spectral densities corresponding to the selected structures are obtained as

\begin{equation}
\rho_1(s, s', q^2) = \int_0^1 dx \int_0^{1-x} dy \left\{ -\frac{3(m_\nu(4 + 6y - 6x) + m_s(-2 + y - x) + 2m_\nu(y - x))}{64\pi^2} \right\}
\end{equation}
where \( M \) in the initial and final mesonic channels, respectively.

Having calculated both the hadronic and OPE sides of the correlation function, we match the coefficients of the selected structures from both sides and apply a double-Borel transformation. As a result, we get the following sum rules for the form factors:

\[
\rho_2(s', q^2) = \int_0^1 dx \int_0^{1-x} dy \left\{ \Theta[y-x](m_B x + m_c (-1 + 2y + x)) \right\} \Theta[L(s', q^2)],
\]

\[
\rho_3(s', q^2) = \int_0^1 dx \int_0^{1-x} dy \left\{ -3(m_B x(-2 + 3y + x) + m_c (2y^2 + y(-1 + x) + (-1 + x)x)) \right\} \times \Theta[L(s', q^2)],
\]

\[
\rho_4(s', q^2) = 0,
\]

(19)

where \( \Theta[... \right\} \Theta[L(s', q^2)], \)

\[
L(s', q^2) = -m_B^2 y - s y(x + y - 1) - x \left( m_B^2 - q^2 y + s(x + y - 1) \right).
\]

(20)

The \( \Pi_{\text{nonpert}}^i(q^2) \) functions are obtained up to five dimension operators. As they have very lengthy expressions, we do not show their explicit forms here.

Having calculated both the hadronic and OPE sides of the correlation function, we match the coefficients of the selected structures from both sides and apply a double-Borel transformation. As a result, we get the following sum rules for the form factors:

\[
K(q^2) = \frac{8(m_B + m_s)}{\Delta} \frac{1}{f_B f_{D_s^*}^2 m_{D_s^*}^2 m_B^2} \frac{m_B^2}{m^2} e^{m_B^2/m^2}
\]

\[
\times \left\{ \int_{(m_B + m_s)^2}^{s_0} ds \int_{(m_c + m_s)^2}^{s_0'} ds' \rho_1(s, s', q^2) + \hat{B}\Pi_{1}^{\text{nonpert}} \right\},
\]

\[
b_-(q^2) = \frac{12(m_B + m_s)}{f_B f_{D_s^*}^2 m_{D_s^*}^2 m_B^2} \frac{m_B^2}{m^2} e^{m_B^2/m^2}
\]

\[
\times \left\{ \int_{(m_B + m_s)^2}^{s_0} ds \int_{(m_c + m_s)^2}^{s_0'} ds' \rho_2(s, s', q^2) + \hat{B}\Pi_{2}^{\text{nonpert}} \right\} + \frac{\Delta}{\Delta'} K(q^2),
\]

\[
b_+(q^2) = \frac{12(m_B + m_s)}{f_B f_{D_s^*}^2 m_{D_s^*}^2 m_B^2} \frac{m_B^2}{m^2} e^{m_B^2/m^2}
\]

\[
\times \left\{ \int_{(m_B + m_s)^2}^{s_0} ds \int_{(m_c + m_s)^2}^{s_0'} ds' \rho_3(s, s', q^2) + \hat{B}\Pi_{3}^{\text{nonpert}} \right\} - \frac{\Delta - 4m_{D_s^*}^2}{\Delta'} K(q^2),
\]

\[
h(q^2) = \frac{8(m_B + m_s)}{\Delta - 4m_{D_s^*}^2} \frac{1}{f_B f_{D_s^*}^2 m_{D_s^*}^2 m_B^2} \frac{m_B^2}{m^2} e^{m_B^2/m^2}
\]

\[
\times \left\{ \int_{(m_B + m_s)^2}^{s_0} ds \int_{(m_c + m_s)^2}^{s_0'} ds' \rho_4(s, s', q^2) + B\Pi_{4}^{\text{nonpert}} \right\},
\]

(21)

where \( M^2 \) and \( M'^2 \) are the Borel mass parameters; and \( s_0 \) and \( s_0' \) are continuum thresholds in the initial and final mesonic channels, respectively.
3 Numerical results

In this section we present our numerical results for the transition form factors derived from
QCD sum rules and search for the behavior of the these quantities in terms of $q^2$. To obtain
numerical values, we use some input parameters presented in table 1.

| Parameters | Values |
|------------|--------|
| $m_s$      | $(95 \pm 5) \text{ MeV}[25]$ |
| $m_b$      | $(4.18 \pm 0.03) \text{ GeV}[25]$ |
| $m_c$      | $(1.275 \pm 0.025) \text{ GeV}[25]$ |
| $m_{B_s}$  | $(5366.77 \pm 0.24) \text{ MeV } [25]$ |
| $m_{D_s^*(2573)}$ | $(2571.9 \pm 0.8) \text{ MeV } [25]$ |
| $f_{B_s}$  | $(222 \pm 12) \text{ MeV } [26]$ |
| $f_{D_s^*(2573)}$ | $(0.023 \pm 0.011) [24]$ |
| $G_F$      | $1.17 \times 10^{-5} \text{ GeV}^{-2}$ |
| $V_{cb}$   | $(41.2 \pm 1.1) \times 10^{-3}$ |
| $\langle 0|\bar{s}s|0\rangle$ | $-(0.8 \pm 0.24)^3 \text{ GeV}^3 [27]$ |
| $m_0^2(1\text{GeV})$ | $(0.8 \pm 0.2) \text{ GeV}^2 [27]$ |
| $\tau_{B_s}$ | $(1.465 \pm 0.031) \times 10^{-12} s [25]$ |

Table 1: Input parameters used in calculations.

To proceed further, we shall find working regions of the four auxiliary parameters,
namely the Borel mass parameters $M^2$ and $M'^2$ and continuum thresholds $s_0$ and $s'_0$, such that the transition form factors weakly depend on these parameters in those regions. The continuum thresholds $s_0$ and $s'_0$ are the energy squares which characterize the beginning of the continuum and depend on the energy of the first excited states in the initial and final channels, respectively. Our numerical calculations point out the following regions for the continuum thresholds $s_0$ and $s'_0$: $29 \text{ GeV}^2 \leq s_0 \leq 35 \text{ GeV}^2$ and $7 \text{ GeV}^2 \leq s'_0 \leq 11 \text{ GeV}^2$.

The working regions for the Borel mass parameters are calculated demanding that both
the higher states and continuum are sufficiently suppressed and the contributions of the
operators with higher dimensions are small. As a result, we find the working regions
10 $\text{ GeV}^2 \leq M^2 \leq 20\text{GeV}^2$ and $5\text{GeV}^2 \leq M'^2 \leq 10\text{GeV}^2$ for Borel mass parameters. We show, as an example, the dependence of the form factor $K(q^2)$ at $q^2 = 1$ on the Borel mass parameters $M^2$ and $M'^2$ in figure 1. With a quick glance to this figure, we see that not only this form factor depicts weak dependence on the Borel parameters on their working regions, but the perturbative contribution constitutes the main part of the total value.

At this stage, we would like to find the behaviors of the considered form factors in terms of $q^2$ using the working regions for the continuum thresholds and Borel mass parameters. Our calculations depict that the form factors are truncated at $q^2 \simeq 5 \text{ GeV}^2$. To extend the results to the whole physical region, we have to find a fit function such that it coincide with the QCD sum rules results at $q^2 = (0 - 5) \text{ GeV}^2$ region. We find that the form factors are
well fitted to the following function:

$$f(q^2) = f_0 \exp \left[ c_1 \frac{q^2}{m_{fit}^2} + c_2 \left( \frac{q^2}{m_{fit}^2} \right)^2 \right], \quad (22)$$
Now we calculate the decay width of the process under consideration. The differential decay width for $B_s \to D_{s2}^{*}(2573)/\bar{\nu}_\tau$ transition is obtained as (see also [28])

$$\frac{d\Gamma}{dq^2} = \frac{\lambda(m_{B_s}, m_{D_{s2}^*}, q^2)}{4m_{D_{s2}^*}^2} \left(\frac{q^2 - m_{B_s}^2}{q^2}\right)^2 \sqrt{\lambda(m_{B_s}, m_{D_{s2}^*}, q^2)G_F V_{cb}^2} \left\{ \frac{1}{2q^2} \left[ 3m_{B_s}^2 \lambda(m_{B_s}, m_{D_{s2}^*}, q^2) |V_0(q^2)|^2 \right. \right.$$

$$\left. + \frac{1}{2m_{D_{s2}^*}} \left[ (m_{B_s} - m_{D_{s2}^*} - q^2)(m_{B_s} - m_{D_{s2}^*}) V_1(q^2) - \frac{\lambda(m_{B_s}, m_{D_{s2}^*}, q^2)}{m_{B_s} - m_{D_{s2}^*}} V_2(q^2) \right] \right| + \frac{2}{3} (m_{B_s}^2 + q^2) \lambda(m_{B_s}, m_{D_{s2}^*}, q^2) \left[ \frac{A(q^2)}{m_{B_s} - m_{D_{s2}^*}} - \frac{(m_{B_s} - m_{D_{s2}^*}) V_1(q^2)^2}{\sqrt{\lambda(m_{B_s}, m_{D_{s2}^*}, q^2)}} \right.$$

Table 2: Parameters appearing in the fit function of the form factors.

|        | $f_0$         | $c_1$        | $c_2$        | $m_{fit}^2$ (GeV$^2$) |
|--------|---------------|--------------|--------------|-----------------------|
| $K(q^2)$ | 0.85 ± 0.25   | 1.06 ± 0.30  | 0.62 ± 0.17  | 28.80 ± 0.01          |
| $b_+(q^2)$ | (0.09 ± 0.03) GeV$^{-2}$ | 1.73 ± 0.51  | 28.90 ± 8.09 | 28.80 ± 0.01          |
| $b_-(q^2)$ | (3.08 ± 0.83) × 10$^{-4}$ GeV$^{-2}$ | 31.53 ± 8.83 | −42.22 ± 12.67 | 28.80 ± 0.01          |
| $h(q^2)$  | (−0.011 ± 0.003) GeV$^{-2}$ | 1.31 ± 0.38  | 0.010 ± 0.003 | 28.80 ± 0.01          |

Figure 1: **Left:** $K(q^2 = 1)$ as a function of the Borel mass $M^2$ at average values of the $s_0$, $s'_0$ and $M^2$. **Right:** $K(q^2 = 1)$ as a function of the Borel mass $M'^2$ at average values of the $s_0$, $s'_0$ and $M^2$. 

where, the values of the parameters $f_0$, $c_1$, $c_2$ and $m_{fit}^2$ are presented in table 2.
Figure 2: $K(q^2)$ as a function of $q^2$ at $M^2 = 15 GeV^2$, $M'^2 = 7.5 GeV^2$, $s_0 = 35 GeV^2$ and $s'_0 = 9 GeV^2$.

$$K(q^2) = A(q^2) + \left| \frac{A(q^2)}{m_{B_s} - m_{D^*_s}} + \frac{(m_{B_s} - m_{D^*_s})V_1(q^2)}{\sqrt{\lambda(m_{B_s}^2, m_{D^*_s}^2, q^2)}} \right|^2, \quad (23)$$

where

$$A(q^2) = -(m_{B_s} - m_{D^*_s})h(q^2),$$
$$V_1(q^2) = -\frac{K(q^2)}{m_{B_s} - m_{D^*_s}},$$
$$V_2(q^2) = (m_{B_s} - m_{D^*_s})b_+(q^2),$$
$$V_0(q^2) = \frac{m_{B_s} - m_{D^*_s}^2}{2m_{D^*_s}}V_1(q^2) - \frac{m_{B_s} + m_{D^*_s}^2}{2m_{D^*_s}}V_2(q^2) - \frac{q^2}{2m_{D^*_s}}b_-(q^2),$$
$$\lambda(a, b, c) = a^2 + b^2 + c^2 - 2ab - 2ac - 2bc. \quad (24)$$

Performing the integral over $q^2$ in the above equation at whole physical region, finally, we obtain the values of the total decay widths and branching ratios for all lepton channels as presented in table 3. The quoted errors in results are due to the errors in determinations of the working regions of the continuum thresholds, Borel mass parameters as well as uncertainties coming from other input parameters. The orders of branching fractions show that the semileptonic $B_s \to D^*_s(2573)\ell \bar{\nu}_\ell$ is accessible, experimentally at all lepton channels in near future.

In summary, we have calculated the transition form factors governing the semileptonic $B_s \to D^*_s(2573)\ell \bar{\nu}_\ell$ transition at all lepton channels using an appreciate three-point correlation function. The fit functions of the form factors have been used to estimate the
Table 3: Numerical results for the decay widths and branching ratios at different lepton channels.

| fit function | $\Gamma$(GeV) | $Br$ |
|--------------|---------------|------|
| $B \to D_{s2}^*(2573)\tau\bar{\nu}_\tau$ | $(3.56 \pm 1.07) \times 10^{-16}$ | $(0.79 \pm 0.24) \times 10^{-3}$ |
| $B \to D_{s2}^*(2573)\mu\bar{\nu}_\mu$ | $(3.13 \pm 0.91) \times 10^{-15}$ | $(0.69 \pm 0.21) \times 10^{-2}$ |
| $B \to D_{s2}^*(2573)e\bar{\nu}_e$ | $(3.15 \pm 0.95) \times 10^{-15}$ | $(0.70 \pm 0.21) \times 10^{-2}$ |

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