$g_{B^*B\rho}$ and $g_{D^*D\rho}$ coupling constants in light cone QCD sum rules

T. M. Aliev,* D. A. Demir, E. Iltan, and N. K. Pak
Physics Department, Middle East Technical University, Ankara, Turkey
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The decay constants for the strong off-shell $B^* \to B\rho$ and $D^* \to D\rho$ decays are calculated in the framework of light cone QCD sum rules. The results are shown to be in agreement with the predictions of the "classical" sum rules method, and with those of the vector dominance model.

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I. INTRODUCTION

Determination of the various characteristics of the heavy flavored hadrons experiments requires information about the physics at large distance. Although the exclusive decays of heavy flavored hadrons are often easier to measure experimentally, for the interpretation of the data one needs accurate estimations of form factors and other matrix elements. There exists a number of nonperturbative approaches for such calculations. Among these, the QCD sum rules [1] are one of the most powerful methods.

The aim of this work is the determination of the $B^*(D^*) \to B(D)\rho$ coupling constant in the framework of QCD sum rules. Here we use an alternative to the "classical" sum rules method, namely the QCD sum rules on light cone [3]. This approach is interesting in several respects: First, all the symmetries of the theory, such as gauge, Lorentz, and conformal invariances [4,5], are preserved in the coordinate space. Second, the covariant light cone expansion in powers of "x^2", which is completely different from the usual Wilson operator product expansion (OPE) based on the T product of currents at small distances, allows us to separate the higher twist effects. It is well known that the OPE on the light cone is performed over the twist of the operators, instead of dimensions, and the main contribution comes from the operators with minimal twist. Matrix elements of nonlocal operators sandwiched between a hadronic state and the vacuum give hadron wave function of increasing twist. The advantage of this approach is that it provides additional information about high-energy asymptotics of correlation functions in QCD, which is accumulated in the wave functions. The high energy behavior of these functions is dictated by the approximate conformal invariance of QCD.

This method was successfully applied for estimating the decay rate of the radiative decay $\Sigma \to p\gamma$ [6], nucleon magnetic moments, the strong couplings $g_{\pi NN}$, $g_{\rho\omega\pi}$ [7], form factors of semileptonic and radiative B-D meson decays [8], the couplings $g_{B^*B\pi}$, $g_{D^*D\pi}$ [9], and the $\pi A\gamma^*$ form factors [10].

In the present paper we calculate the strong coupling constants, $g_{B^*B\rho}$ and $g_{D^*D\rho}$, using the light cone sum rule. The paper is organized as follows: In Sec. II we derive the light cone sum rule for the $B^*(D^*) \to B(D)\rho$ coupling constants. Section III is devoted to the analysis of the sum rules and discussions.

II. CALCULATION OF THE $g_{B^*B\rho}$ AND $g_{D^*D\rho}$ COUPLING CONSTANTS

According to the QCD sum rule ideology, for the calculation of the $g_{B^*B\rho}$ coupling constant it is necessary to construct a suitable correlation function in hadronic and quark-gluon languages. For this purpose, let us consider the following correlator function information about high-energy asymptotics of correlation function

$$F_\mu = i\int d^4x e^{ipx} \times\langle\rho(q,\epsilon)|T\{\bar\psi(x)\gamma_\mu b(x), \bar b(0)i\gamma_5\psi(0)\}|0\rangle. \quad (1)$$

Here ψ and b are the light and beauty quark fields.

When the ρ meson is on the mass shell, $q^2 = m_\rho^2$, the correlation function (1) depends on two variables, p^2 and $(p+q)^2$.

The correlator (1) in quark language can be calculated in the deep Euclidian region where both variables p^2 and $(p+q)^2$ are negative and large, so that the heavy quark is sufficiently far off-shell. Therefore we can use the perturbative expansion of its propagator in the external field with slowly varying fluctuations inside the ρ meson. The leading contribution is represented by the diagram in Fig. 1. For the calculation of this diagram we use the free-heavy quark propagator:

$$\langle 0|T\{\bar b(x)b(0)\}|0\rangle = iS_b^0(x) = \int\frac{d^4k}{(2\pi)^4 i}e^{-ikx}\frac{\not k + m_b}{m_b^2 - k^2}. \quad (2)$$

Substituting Eq. (2) in Eq. (1), for the leading contribution, we get

*Permanent address: Institute of Physics, Azerbaijanian Academy of Sciences, Baku, Azerbaijan.

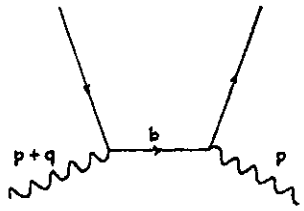

FIG. 1. Diagrams contributing to the correlation function Eq. (1). Solid lines represent quarks, wavy lines external currents.

$$F_\mu = i\int \frac{d^4k d^4x}{(2\pi)^4 i}\frac{e^{i(p-k)x}}{m_b^2-k^2} \times\langle\rho(q,\epsilon)|\bar{u}(x)\gamma_\mu(\not{k}+m_b)i\gamma_5 u(0)|0\rangle. \quad (3)$$

Thus, in general we deal with matrix elements of gauge-invariant nonlocal operators, sandwiched between the vacuum and the meson states. These matrix elements define meson wave functions on the light cone. From Eq. (3) it follows that we need two matrix elements:

$$\langle\rho(q,\epsilon)|\bar{u}(x)\gamma_\mu\gamma_\alpha i\gamma_5 u(0)|0\rangle \quad (4)$$

and

$$\langle\rho(q,\epsilon)|\bar{u}(x)\gamma_\mu i\gamma_5 u(0)|0\rangle. \quad (5)$$

For the calculation of the matrix element (4), we use the Fierz identity

$$\gamma_\mu\gamma_\alpha\gamma_5 = -\epsilon_{\mu\alpha\lambda\rho}\sigma_{\lambda\rho} + g_{\mu\alpha}\gamma_5$$

and get

$$\langle\rho(q,\epsilon)|\bar{u}(x)\gamma_\mu\gamma_\alpha\gamma_5 u(0)|0\rangle = -\epsilon_{\mu\alpha\lambda\rho}\langle\rho(q,\epsilon)|\bar{u}(x)\sigma_{\lambda\rho}u(0)|0\rangle + g_{\mu\alpha}\langle\rho(q,\epsilon)|\bar{u}(x)\gamma_5 u(0)|0\rangle. \quad (6)$$

By definition, the matrix element proportional to $g_{\mu\alpha}$ in Eq. (6) is equal to zero. The first term in Eq. (6), following [2], can be defined as

$$\langle\rho(q,\epsilon)|\bar{u}(x)\sigma_{\lambda\rho}u(0)|0\rangle = i(e_\lambda q_\rho - e_\rho q_\lambda)f_\rho^\perp \int_0^1 du e^{-iupx}\phi_\perp(u,\mu^2). \quad (7)$$

The matrix element (5) can be defined by the following expression [11]:

$$\langle\rho(q,\epsilon)|\bar{u}(x)\gamma_\mu\gamma_5 u(0)|0\rangle = (1/4)\epsilon_{\mu\alpha\lambda\tau}e_\alpha q_\rho x_\tau f_\rho m_\rho \int_0^1 du e^{-iuqx} g_\perp(u,\mu^2). \quad (8)$$

The function $\phi_\perp(u,\mu^2)$, whose leading twist is two, describes the distribution of the fraction of the total momentum carried by the quark in the transversely polarized ρ meson. In [11], it is shown that the function $g_\perp(u,\mu^2)$ contains the contributions coming from the operators of twist two and three. The twist-three operator contributions are due to the gluon exchange in Fig. 1, and it is described by three-particle quark-antiquark-gluon wave functions of the transversely polarized vector mesons (see, for example, [2,11]). In this work, we do not consider these contributions and take into account only the leading twist operators. In this accuracy, we take $g_\perp(u) = 6u(1-u)$ [11], in our numerical calculations.

In [11] it is shown that the contributions to the transversal $g_\perp$ can be expressed in terms of the longitidunal wave function $\phi_{||}$ with leading twist $\tau = 2$ defined as

$$\langle 0|\bar{\psi}(0)\gamma_\mu\psi(x)|\rho(\lambda,p)\rangle = p_\mu\frac{e^\lambda x}{(p\cdot x)}f_\rho m_\rho\int_0^1 du e^{-iupx}\phi_{||}(u,\mu^2).$$

Namely,

$$\frac{d}{du}g_\perp^{\tau=2}(u) = 2\left[\int_0^u dv\frac{\phi_{||}(v)}{\bar{v}} + \int_u^1 dv\frac{\phi_{||}(v)}{v}\right],$$

where $\bar{v} = 1-v$.

In (7) and (8), u is the difference in the fractions of the ρ meson momentum carried by the quark and antiquark, and μ is the renormalization point of the wave function.

Using Eqs. (3), (7), and (8) we get for F_μ

$$F_\mu(p,q) = \int_0^1 du\int\frac{d^4x d^4k}{(2\pi)^4}e^{i(p-k+uq)x}\frac{i}{m_b^2-k^2} \times\{m_b[-\tfrac{1}{4}\epsilon_{\mu\alpha\rho\tau}e_\alpha q_\rho x_\tau f_\rho m_\rho g_\perp(u,\mu^2)] + \epsilon_{\mu\alpha\lambda\rho}ie_\lambda q_\rho k_\alpha f_\rho^\perp\phi_\perp(u,\mu^2)\}. \quad (9)$$

Writing x_β in the momentum space as $x_\beta = -i\frac{\partial}{\partial p_\beta}$, and performing the integrations over x and k variables, we get

$$F_\mu = \epsilon_{\mu\alpha\beta\sigma}q_\beta e_\alpha p_\sigma F, \quad (10)$$

where F is an invariant function of variables p^2 and $(p+q)^2$, and has the following explicit form:

$$F(p^2,(p+qu)^2) = \int_0^1 du\left\{\frac{1}{2}\frac{m_b f_\rho m_\rho g_\perp(u,\mu^2)}{[m_b^2-(p+qu)]^2} + 2f_\rho^\perp\phi_\perp(u,\mu^2)\frac{1}{m_b^2-(p+qu)^2}\right\}. \quad (11)$$

Now we apply the double Borel transformation

$$\hat{B} = \lim_{\substack{n,k\to\infty,\\ -p^2\to\infty,\,-p'^2\to\infty,\\ -\frac{p^2}{n}=M^2,\,-\frac{p'^2}{k}=M'^2}} \frac{1}{(n-1)!(k-1)!}(-p^2)^n\left(\frac{d}{dp^2}\right)n(-p'^2)^k\left(-\frac{d}{dp'^2}\right)k \tag{12}$$

to Eq. (11). For this purpose we use the exponential representation:

$$\frac{1}{A^n} = \frac{1}{(n-1)!}\int_0^\infty \alpha^{n-1}e^{-\alpha A}d\alpha. \tag{13}$$

Here we would like to present the technical details of our calculation. Let us first consider the first term in Eq. (11):

$$F_1 = \frac{1}{2}\int_0^1 du \frac{m_b f_\rho m_\rho g_\perp(u,\mu^2)}{[m_b^2-(p+qu)^2]^2}. \tag{14}$$

In the denominator, going into the Eucledian region, we note

$$(p+uq)^2 = (1-u)p^2 + u(p+q)^2 + m_\rho^2 u(1-u).$$

Then using Eq. (13), we get the expression

$$F_1 = \frac{1}{2}\int_0^1 du\int_0^\infty d\alpha\alpha e^{-\alpha[m_b^2+p^2(1-u)+(p+q)^2u+m_\rho^2u(1-u)]}m_b f_\rho m_\rho g_\perp g(u,\mu^2). \tag{15}$$

Applying the double Borel transformation and using

$$B_{M^2}(p^2)e^{-\alpha p^2(1-u)} = \delta[1-\alpha M^2(1-u)]$$

(for more detail see [12]) we obtain

$$F_1 = \frac{1}{2}\int_0^1 du\int_0^\infty \alpha d\alpha\delta[1-\alpha M^2(1-u)]\delta(1-\alpha M'^2u)e^{-\alpha[m_b^2+m_\rho^2u(1-u)]}m_b f_\rho m_\rho g_\perp g(u,\mu^2).$$

After performing integrations over α and u we get the following expression:

$$F_1 = \frac{1}{2}m_b f_\rho m_\rho g_\perp g(u,\mu^2)\frac{1}{M^2M'^2}e^{-\frac{1}{M'^2u}[m_b^2+m_\rho^2u(1-u)]}\Big|_{u=\frac{M^2}{M'^2+M^2}}.$$

For the second term in Eq. (11), we follow the same procedure and we get the final result for the invariant function F^{theor}:

$$F^{\text{theor}} = \frac{1}{M^2M'^2}e^{-\frac{1}{M'^2u}[m_b^2+m_\rho^2u(1-u)]}\left[\frac{1}{2}m_b f_\rho m_\rho g_\perp(u,\mu^2) + 2\phi_\perp(u,\mu^2)f_\rho^\perp M'^2u\right]\Bigg|_{u=\frac{M^2}{M'^2+M^2}}. \tag{16}$$

Since the mass of B^* and B mesons are practically equal, we can take $M^2 = M'^2 = 2M_*^2$ (see also [11]). In this case $u = 1/2$ and the invariant function F^{theor} becomes

$$F^{\text{theor}}(M_*^2) = \frac{e^{-\frac{1}{M_*^2}(m_b^2+m_\rho^2/4)}}{4M_*^2}\left\{\frac{1}{2}m_b f_\rho m_\rho g_\perp g(1/2,\mu^2)/M_*^2 + 2f_\rho^\perp\phi_\perp(1/2,\mu^2)\right\}.$$

At this point we subtract the continuum contribution by replacing the exponential factor $e^{-(m_b^2+m_\rho^2/4)/M_*^2}$ by $e^{-(m_b^2/M_*^2+m_\rho^2/4)} - e^{-s_0/M_*^2}$ [9] where s_0 is the continuum threshold in the B channel, $s_0 = 36$ GeV^2. So, for the theoretical part we get

$$F(M_*^2) = \frac{1}{4M_*^2}\left[e^{-\frac{1}{M_*^2}(m_b^2+m_\rho^2/4)} - e^{-s_0/M_*^2}\right]\left\{\frac{1}{2}m_b f_\rho m_\rho g_\perp g(1/2,\mu^2)/M_*^2 + 2f_\rho^\perp\phi_\perp(1/2,\mu^2)\right\}. \tag{17}$$

Note that $u = 1/2$ corresponds to the case where the probabilities of the fractions of the ρ meson momentum carried by quark and antiquark are equal.

Now we need the physical part of the sum rules. Saturating Eq. (1) by B^* and B mesons we get the expression for the physical part,

$$F_\mu(p,q) = \epsilon_{\mu\alpha\beta\sigma} q_\alpha e^\rho_\beta p_\sigma \frac{m_{B^*} f_{B^*} f_B m_B^2}{m_b} \frac{g_{B^*B\rho}}{(p^2 - m_{B^*}^2)[(p+q)^2 - m_B^2]}. \quad (18)$$

Here q and ϵ are the four-momentum and the polarization vector of the ρ meson; p is the four-momentum of $B^*(D^*)$ meson. In deriving Eq. (18), we use the following definitions:

$$\langle B^*|B\rho\rangle = g_{B^*B\rho}\epsilon_{\lambda\alpha\beta\sigma} q_\alpha e^\rho_\beta p_\sigma e^{B^*}_\lambda ,$$
$$\langle 0|\bar\psi\gamma_\mu b|B^*\rangle = f_{B^*} m_{B^*} e^{B^*}_\mu ,$$
$$\langle B|\bar b i\gamma_5 \psi|0\rangle = \frac{m_B^2 f_B}{m_b}. \quad (19)$$

Applying the double Borel transformation to Eq. (18), and equating the physical invariant function to the theoretical one, Eq. (17), we obtain the sum rule for the coupling constant $g_{B^*B\rho}$:

$$g_{B^*B\rho} f_{B^*} f_B = \frac{4M_*^4 m_b}{m_{B^*} m_B^2} \exp(m_{B^*}^2 + m_B^2/2M_*^2) F^{\text{theor}}. \quad (20)$$

In numerical calculations we use the following expressions for the ρ meson wave functions [11]:

$$\phi_\perp(u,\mu^2) = 6u(1-u)[1 + a_2(\mu)(\xi^2 - 1/5) + \cdots], \quad (21)$$
$$g_\perp(u,\mu^2) = 6u(1-u), \quad (22)$$

where

$$a_2(\mu) = a_2(\mu_0)\left(\frac{\alpha_s(\mu^2)}{\alpha_s(\mu_0^2)}\right)^{\gamma_2/\beta},$$

and $\xi = 2u - 1$, $\beta = \frac{1}{3}(11N_c - 2n_f)$ and the anomalous dimension γ_2 is given by [13]

$$\gamma_n = \frac{N_c^2 - 1}{2N}\left[1 + 4\sum_{j=2}^{n+1}\frac{1}{j}\right].$$

The coefficients in the Gegenbauer expansion Eq. (21) at low scale, $a_n(\mu_0)$, should be determined by a certain nonperturbative method, or should be extracted from the experimental data. We use the model given in [2] for ρ meson wave functions. At $\mu_0^2 = 1$ GeV^2 this model predicts

$$a_2(\mu_0^2 = 1\ \text{GeV}^2) = -1.25.$$

Using this value, at $\mu^2 = m_b^2$ scale we find that $a_2(\mu^2 = m_b^2) = -0.85$ at $\Lambda^{(5)} = 225$ MeV. For the determination of the coupling constant $g_{D^*D\rho}$, we make the following replacements: $b \to c$, $B^* \to D^*$, $B \to D$, and $a_2(m_b^2) \to a_2(m_c^2)$.

III. NUMERICAL ANALYSIS AND DISCUSSION

The aim of the present work is to determine the coupling constants $g_{B^*B\rho}$ and $g_{D^*D\rho}$. In the numerical analysis we have used the following input parameters:

$$m_b = 4.6\text{–}4.8 \text{ GeV } [14,15],$$
$$m_c = 1.3\text{–}1.4 \text{ GeV } [1,14],$$
$$m_{B^*(D^*)} = 5.324\,(2.10) \text{ GeV},$$
$$m_{B(D)} = 5.278\,(1.864) \text{ GeV } [16],$$
$$f_\rho^\perp = 0.2\ [2], \quad f_\rho = m_\rho/(9.44\pi)^{1/2}\ [1].$$

According to the QCD sum rules method we have to find a region of M_*^2 where $g_{B^*B\rho} f_{B^*} f_B$ does not practically depend on M_*^2, and at the same time the continuum contribution remains under control, i.e., it constitutes about 30–40 % of the bare loop contribution.

The M_*^2 dependence of $g_{B^*B\rho} f_{B^*} f_B$ at $s_0 \sim 36$ GeV is presented in Fig. 2. The best stability region is 10 GeV$^2 \le M_*^2 \le 20$ GeV^2, and the prediction is

$$g_{B^*B\rho} f_{B^*} f_B = 0.33. \quad (23)$$

Performing the similar calculations for $g_{D^*D\rho} f_{D^*} f_D$ we get (Fig. 3)

$$g_{D^*D\rho} f_{D^*} f_D = 0.57 \quad (24)$$

in the stability region 4 GeV$^2 \le M_*^2 \le 6$ GeV^2, and at $s_0 = 6$ GeV^2. Now let us compare the results (21) and (22) with the predictions on these coupling constants in the framework of "classical sum rules method." In [17],

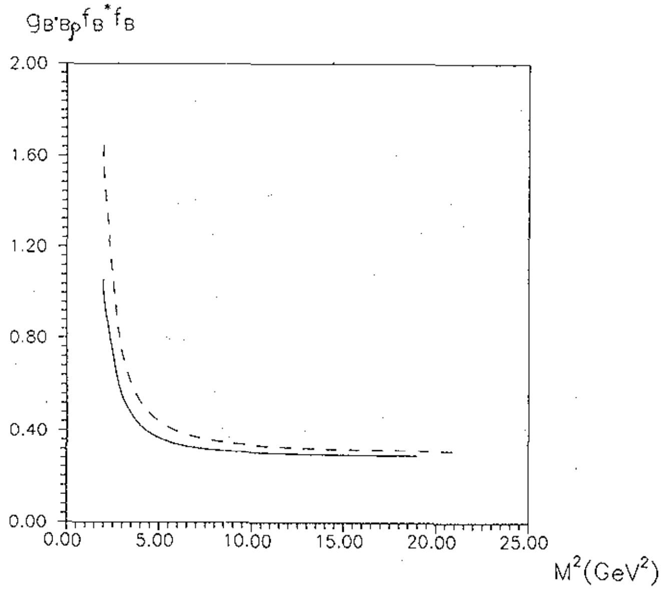

FIG. 2. The dependence of the coupling constant $g_{B^*B\rho} f_{B^*} f_B$ on the Borel parameter square M^2. Solid line corresponds to first set and dashed line to second set of values of the leptonic decay constants.

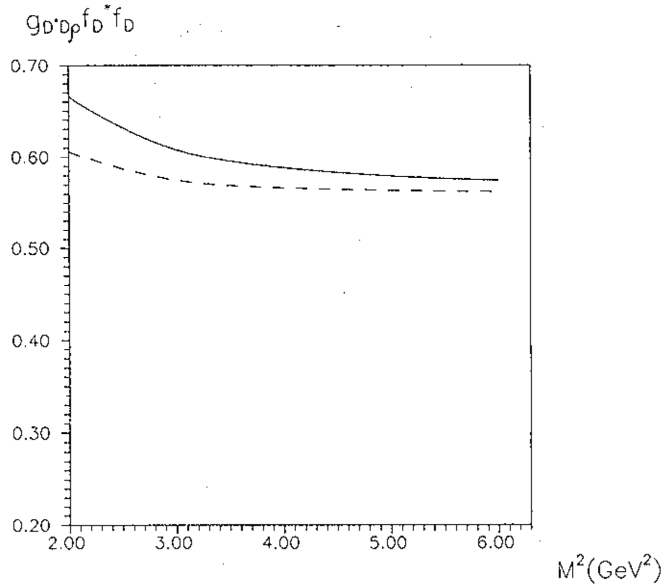

FIG. 3. The same as in Fig. 2 but for the $g_{D^*D\rho}f_{D^*}f_D$ case.

we find that

$$g_{B^*B\rho}f_{B^*}f_B = 0.40,$$
$$g_{D^*D\rho}f_{D^*}f_D = 0.41. \quad (25)$$

We see that the results obtained in these two approaches (the light cone and the classical sum rule) are close. The differences are due to the higher twist effects.

For determination of the coupling constants from Eqs. (23) and (24), we use the following two sets of values for the leptonic decay constants:

$$(1)\ f_B = 140 \text{ MeV } [9,18],\quad f_{B^*} = 160 \text{ MeV } [18],\quad f_D = 170 \pm 10 \text{ MeV } [9,18],\quad f_{D^*} = 240 \pm 20 \text{ MeV } [9], \quad (26)$$

$$(2)\ f_B = 198 \pm 15 \text{ MeV } [19,20,21],\quad f_{B^*} = 213 \text{ MeV } [19],\quad f_D = 180 \text{ MeV } [19,20,21],\quad f_{D^*} = 258 \text{ MeV } [19]. \quad (27)$$

Using Eqs. (23) and (24) and the values of leptonic decay constants we obtain

$$(1)\ g_{B^*B\rho} = 15,\quad g_{D^*D\rho} = 14,$$
$$(2)\ g_{B^*B\rho} = 9,\quad g_{D^*D\rho} = 12.$$

For the calculations of the decay constants in the first set, we used the values of c- and b-quark masses as $m_c = 1.3$ GeV and $m_b = 4.8$ GeV, and $m_c = 1.42$ GeV and $m_b = 4.6$ GeV for the second set.

Finally, we would like to compare our results with the predictions of the vector dominance model (VDM). For the $g_{B^*B\rho}$ and $g_{D^*D\rho}$ coupling constants, VDM predicts a value $g_{B^*B\rho} = g_{D^*D\rho} = \frac{2}{f_\pi} = 16$ [22], where $f_\pi = 133$ MeV is the pion decay constant. We see that if we choose the first set of values of the leptonic decay constants f_B and f_D, our predictions for $g_{B^*B\rho}$ and $g_{D^*D\rho}$ coupling constants are in very good agreement with the VDM predictions.

In conclusion, the coupling constants $g_{B^*B\omega}$, $g_{B^*BK^*}$, $g_{D^*DK^*}$, and $g_{D^*D\omega}$ can be easily obtained by using the similar calculations.

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