Width of the exotic $X_b(5568)$ state through its strong decay to $B_s^0\pi^+$

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The width of the newly observed exotic state $X_b(5568)$ is calculated via its dominant strong decay to $B_s^0\pi^+$ using the QCD sum rule method on the light-cone in conjunction with the soft-meson approximation. To this end, the vertex $X_bB_s\pi$ is studied and the strong coupling $g_{X_bB_s\pi}$ is computed employing for $X_b(5568)$ state the interpolating diquark-antidiquark current of the $[su][\overline{bd}]$ type. The obtained prediction for the decay width of $X_b(5568)$ is confronted and a nice agreement found with the experimental data of the D0 Collaboration.

I. INTRODUCTION

Starting from the discovery of the charmoniumlike resonance $X(3872)$ by Belle Collaboration [1], later confirmed by some other experiments [2–4], investigation of the exotic states became one of the interesting directions in hadron physics. The exotic states, that is particles inner structure of which cannot be described within the two- or three quark scheme of standard hadron spectroscopy, provide rich information to check contemporary theories claiming to explain variety of new phenomena. They can be produced in numerous exclusive and inclusive processes icluding the B meson decays and $\overline{p}p$ collisions. Now, the hadron spectroscopy contains the numerous family of the exotic states XYZ, discovered and studied during past years. Relevant experimental investigations are concentrated on measurements of the masses and decay widths of these states, on exploration of their spins and parities. At the same time, theoretical studies are aimed to invent new approaches and methods to calculate parameters of the exotic particles (see for instance [5–12] and references therein).

There are various models in the literature suggested to reveal the internal quark-antiquark structure of the exotic states and explain the wide range of corresponding experimental data. One of the popular models is the four-quark or tetraquark picture of the exotic particles. In accordance with this approach, new charmoniumlike states are composed of two heavy and two light quarks. These quarks may cluster into the colored diquark and antidiquark, which are organized in such a way that to reproduce quantum numbers of the corresponding exotic states. Other possibilities of the organization of the exotic particles in the context of the tetraquark picture include a meson-molecule and hadro-quarkonium models. Among alternative views on the nature of the exotic states one should mention the conventional charmonium model and its hybrid extensions. All of these models consider the exotic states as particles containing a $c\bar{c}$ component. In fact, the known exotic states consist of $c\bar{c}$ and two light quarks. In other words, the number of the quark flavors inside of the known four-quark exotic particles does not exceed three.

Recently, the D0 Collaboration in Ref. [13] reported the observation of a narrow structure $X(5568)$ in the decays $X(5568) \to B_s^0\pi^{\pm}$, $B_s^0 \to J/\psi\phi$, $J/\psi \to \mu^+\mu^-$, $\phi \to K^+K^-$. This result was based on $p\bar{p}$ data collected at the Fermilab Tevatron at $\sqrt{s} = 1.96$ TeV. The mass of the new state extracted from the experiment equals to $m_{X_b} = 5567.8 \pm 2.9(\text{stat})^{+0.9}_{-1.9}(\text{syst})$ MeV, whereas its decay width was estimated as $\Gamma = 21.9 \pm 6.4(\text{stat})^{+5.0}_{-2.5}(\text{syst})$ MeV. The exotic state $X(5568)$ is supposedly a scalar particle with the quantum numbers $J^{PC} = 0^{++}$ and built of four distinct quark flavors. In fact, from the existing decay channel $X(5568) \to B_s^0\pi^{\pm}$ one can conclude that the state $X(5568)$ contains the valence b, s, u and d quarks. This state can be described as the quark-antiquark bound state with one of the possible structures $[bu][\overline{ds}]$, $[bd][\overline{su}]$, $[su][\overline{bd}]$ and $[sd][\overline{bu}]$, or may be considered as a molecule composed of B and $\bar{K}$ mesons [13].

To differ the $X(5568)$ from the conventional members of the X family of exotic particles, in Ref. [14] we introduced the notation $X_b(5568)$. In the present work we will use this abbreviation, as well. In Ref. [14], adopting the diquark-antidiquark structure $X_b = [su][\overline{bd}]$ we calculated, for the first time, the mass and decay constant of the $X_b(5568)$ state. We employed QCD two-point sum rule method and taken into account the vacuum condensates up to eight dimensions. Our prediction for the mass of the $X_b(5568)$ state is in a nice agreement with data of the D0 Collaboration.

The mass and pole residue of the exotic state $X(5568)$ were also calculated in Ref. [15]. The mass of the state $X(5568)$ was evaluated in Refs. [16–18], as well. To perform relevant calculations, in these works different forms of diquark-antidiquark interpolating currents were used. The obtained results for m_X agree with each other, and are consistent with experimental data of the D0 Collaboration.

In the present paper we extend our investigation of the newly observed state $X_b(5568)$ by calculating the width of the dominant decay $X_b(5568) \to B_s^0\pi^+$. To this

end, we compute the strong coupling $g_{X_b B_s \pi}$ by applying methods of QCD light-cone sum rule (LCSR) and soft-meson approximation [19–21]. The soft-meson approximation is required because the X_b state contains the four valence quarks, and as a result, the light-cone expansion of the correlation functions reduces to the short-distance expansion in terms of local matrix elements. This approximation was applied in our previous work Ref. [22] to calculate the decay widths of the $Z_c(3900)$ state, where a good agreement with the experimental data and available theoretical results was found.

This paper is structured in the following way. In section II, we calculate the strong coupling $g_{X_b B_s \pi}$ and width of the decay $X_b(5568) \to B_s^0 \pi^{\pm}$. Section III contains our numerical result. Here we compare our prediction for the width of the $X_b(5568)$ state with the relevant experimental data of the D0 Collaboration. This section contains also our concluding remarks.

II. THE VERTEX $X_b B_s \pi$ AND STRONG DECAY $X_b \to B_s^0 \pi^+$

In this section we calculate the width of the $X_b \to B_s^0 \pi$ decay. To this purpose, as the first step we calculate the strong coupling $g_{X_b B_s \pi}$ by means of the QCD light-cone sum rules method applying the soft-meson approximation. In order to get the sum rule expression for the coupling $g_{X_b B_s \pi}$ we consider the correlation function

$$\Pi(p,q) = i \int d^4 x e^{ipx} \langle \pi(q) | \mathcal{T} \{ J^{B_s}(x) J^{X_b \dagger}(0) \} | 0 \rangle, \quad (1)$$

where the interpolating currents are given as

$$J^{X_b}(x) = \varepsilon^{ijk} \varepsilon^{imn} \left[s^j(x) C \gamma_\mu u^k(x) \right] \left[\overline{b}^m(x) \gamma^\mu C \overline{d}^n(x) \right], \quad (2)$$

and

$$J^{B_s}(x) = \overline{b}_l(x) i \gamma_5 s_l(x). \quad (3)$$

In Eqs. (2) and (3) i, j, k, m, n and l are the color indices and C is the charge conjugation matrix.

In terms of the physical degrees of freedom the correlation function $\Pi(p,q)$ is determined by the expression

$$\Pi^{\text{Phys}}(p,q) = \frac{\langle 0 | J^{B_s} | B_s(p) \rangle}{p^2 - m_{B_s}^2} \langle B_s^0(p) \pi(q) | X_b(p') \rangle$$
$$\times \frac{\langle X_b(p') | J^{X_b \dagger} | 0 \rangle}{p'^2 - m_{X_b}^2} + \dots, \quad (4)$$

where by dots we denote contributions of the higher resonances and continuum states. Here p, q and $p' = p + q$, are the momenta of B_s^0, π, and the X_b states, respectively. To compute the correlation function we introduce also the matrix elements

$$\langle 0 | J^{B_s} | B_s(p) \rangle = \frac{f_{B_s} m_{B_s}^2}{m_b + m_s},$$
$$\langle X_b(p') | J^{X_b \dagger} | 0 \rangle = f_{X_b} m_{X_b},$$
$$\langle B_s^0(p) \pi(q) | X_b(p') \rangle = g_{X_b B_s \pi} p \cdot p'. \quad (5)$$

In Eq. (5) by m_{X_b} and f_{X_b} we denote the mass and decay constant of the X_b state, whereas m_{B_s} and f_{B_s} are the same parameters of the B_s^0 meson.

Using these matrix elements for the correlation function we obtain

$$\Pi^{\text{Phys}}(p,q) = \frac{f_{B_s} f_{X_b} m_{X_b} m_{B_s}^2 g_{X_b B_s \pi}}{\left(p'^2 - m_{X_b}^2 \right) \left(p^2 - m_{B_s}^2 \right) (m_b + m_s)} p \cdot p'. \quad (6)$$

In the soft-meson limit accepted in the present work $q = 0$, and as a result, $p = p'$. The reason why we apply the soft-meson limit was explained in rather detailed form in our previous article [22]. Nevertheless, for completeness we provide briefly corresponding arguments. In fact, the X_b state and interpolating current Eq. (2) contain four quark fields at the same space-time location. Substitution of this current into the correlation function and subsequent contraction of the b and s quark fields yield expressions, where the remaining light quarks are placed between the π meson and vacuum states forming local matrix elements. Stated differently, we appear in the situation when dependence of the correlation function on the meson distribution amplitudes disappears and integrals over the meson DAs reduce to overall normalization factors. In the context of the QCD LCSR method such situation is possible in the kinematical limit $q \to 0$, when the light-cone expansion is replaced by the short-distant one. As a result, instead of the expansion in terms of DAs, one gets expansion over the local matrix elements [21]. In this limit the relevant invariant amplitudes in the correlation function depend only on the variable p^2.

In the case under consideration the corresponding invariant amplitude reads

$$\Pi^{\text{Phys}}(p^2) = \frac{f_{B_s} f_{X_b} m_{X_b} m_{B_s}^2 g_{X_b B_s \pi}}{\left(p^2 - m_{X_b}^2 \right) \left(p^2 - m_{B_s}^2 \right) (m_b + m_s)} m^2$$
$$+ \Pi^{(\text{RS:C})}(p^2), \quad (7)$$

where $m^2 = \left(m_{X_b}^2 + m_{B_s}^2 \right) / 2$. In Eq. (7), $\Pi^{(\text{RS:C})}(p^2)$ is the contribution arising from the higher resonances and continuum states.

What is also important, instead of the two-variable Borel transformation on p^2 and p'^2, now we have to use the one-variable Borel transformation on p^2: this fact plays a crucial role in deriving sum rules for the strong couplings. Indeed, the soft-meson approximation considerably simplifies the QCD side of the sum rules, but leads to more complicated expression for its hadronic representation. In the soft limit, the ground state contribution can be written in the form

$$\Pi^{\text{Phys}}(p^2) \cong \frac{f_{B_s} f_{X_b} m_{X_b} m_{B_s}^2 g_{X_b B_s \pi}}{\left(p^2 - m^2 \right)^2 (m_b + m_s)} m^2. \quad (8)$$

The Borel transformation on the variable p^2 applied to

this correlation function yields

$$\Pi^{\mathrm{Phys}}(M^2) = \frac{f_{B_s} f_{X_b} m_{X_b} m_{B_s}^2 g_{X_b B_s \pi}}{(m_b + m_s)} m^2$$

$$\times \frac{1}{M^2} e^{-m^2/M^2}. \tag{9}$$

In the soft-meson limit we have to use the one-variable Borel transformation, therefore transitions from the exited states $m^* > m_{B_s}$ to the ground state in the B_s^0 channel (similar arguments are valid for the X_b channel, as well) contribute to the hadronic part of the sum rules. The relevant contributions even after the Borel transformation are not suppressed relative to the ground state one [20, 21]. To remove from the sum rules unsuppressed contributions we employ a prescription elaborated in Ref. [20] and apply the operator

$$\left(1 - M^2 \frac{d}{dM^2}\right) M^2 e^{m^2/M^2} \tag{10}$$

to both sides of the sum rule expression.

To find the QCD side of the sum rules one should calculate the correlation function in the quark-gluon degrees of freedom. Contraction of the s and b-quark fields results in the expression

$$\Pi^{\mathrm{QCD}}(p,q) = -\int d^4x e^{ipx} \varepsilon^{ijk} \varepsilon^{imn} \left[\gamma^\mu \widetilde{S}_s^{lj}(x)\gamma_5\right.$$

$$\left.\times \widetilde{S}_b^{ml}(-x)\gamma_\mu\right]_{\alpha\beta} \langle\pi(q)|\overline{u}_\alpha^k(0)d_\beta^n(0)|0\rangle, \tag{11}$$

where α and β are the spinor indices. In Eq. (11) we introduce the notation

$$\widetilde{S}_{b(s)}^{ij}(x) = C S_{b(s)}^{ijT}(x) C,$$

with $S_s^{ij}(x)$ and $S_b^{ij}(x)$ being the s- and b-quark propagators, respectively.

In general, for calculation of the $\Pi^{\mathrm{QCD}}(p,q)$ we have to use the light-cone expansion for the s- and b-quark propagators. But because in the matrix elements the light quark fields are already fixed at the point $x = 0$, it is enough in calculations to utilize the local propagators. We choose the s-quark propagator $S_q^{ij}(x)$ in the x-space in the form

$$S_q^{ij}(x) = i\delta_{ij} \frac{\slashed{x}}{2\pi^2 x^4} - \delta_{ij} \frac{m_s}{4\pi^2 x^2} - \delta_{ij} \frac{\langle \overline{s}s\rangle}{12}$$

$$+ i\delta_{ij} \frac{\slashed{x} m_s \langle\overline{s}s\rangle}{48} - \delta_{ij} \frac{x^2}{192}\langle\overline{s}g\sigma Gs\rangle + i\delta_{ij}\frac{x^2 \slashed{x} m_s}{1152}\langle\overline{s}g\sigma Gs\rangle$$

$$- i\frac{gG_{ij}^{\alpha\beta}}{32\pi^2 x^2}[\slashed{x}\sigma_{\alpha\beta} + \sigma_{\alpha\beta}\slashed{x}] + \dots \tag{12}$$

For the b-quark propagator $S_b^{ij}(x)$ we employ the expression [23]

$$S_b^{ij}(x) = i\int \frac{d^4k}{(2\pi)^4} e^{-ikx} \left[\frac{\delta_{ij}(\slashed{k}+m_b)}{k^2-m_b^2}\right.$$

$$\left.- \frac{gG_{ij}^{\alpha\beta}}{4}\frac{\sigma_{\alpha\beta}(\slashed{k}+m_b)+(\slashed{k}+m_b)\sigma_{\alpha\beta}}{(k^2-m_b^2)^2}\right] + \dots \tag{13}$$

In Eqs. (12) and (13) the short-hand notation

$$G_{ij}^{\alpha\beta} \equiv G_A^{\alpha\beta} t_{ij}^A, \quad A = 1, 2 \dots 8,$$

is used, where i, j are color indices, and $t^A = \lambda^A/2$ with λ^A being the standard Gell-Mann matrices. In the nonperturbative terms the gluon field strength tensor $G_{\alpha\beta}^A \equiv G_{\alpha\beta}^A(0)$ is fixed at $x = 0$.

To proceed we use the expansion

$$\overline{u}_\alpha^k d_\beta^m \to \frac{1}{4}\Gamma_{\beta\alpha}^j \left(\overline{u}^k \Gamma^j d^m\right), \tag{14}$$

where Γ^j is the full set of Dirac matrixes

$$\Gamma^j = \mathbf{1},\ \gamma_5,\ \gamma_\lambda,\ i\gamma_5\gamma_\lambda,\ \sigma_{\lambda\rho}/\sqrt{2}.$$

In order to fix the local matrix elements necessary in our calculations, first we consider the perturbative component of the b-quark ($\sim \delta_{ml}$) and terms $\sim \delta_{lj}$ from the s-quark propagators and perform the summation over the color indices. To this end, we use the overall color factor $\varepsilon^{ijk}\varepsilon^{imn}$, color factors of the propagators, and the projector onto a color-singlet state $\delta^{km}/3$. As a result, we find that for such terms the replacement

$$\frac{1}{4}\Gamma_{\beta\alpha}^j \left(\overline{u}^k \Gamma^j d^m\right) \to \frac{1}{2}\Gamma_{\beta\alpha}^j \left(\overline{u}\Gamma^j d\right) \tag{15}$$

has to be implemented. In the case of the nonperturbative contributions, which appear as a product of the perturbative part of one propagator with the $\sim G$ component of the another propagator, for example, we obtain

$$\varepsilon^{ijk}\varepsilon^{imn}\delta_{lj}G_{ml}^{\rho\delta}\frac{1}{4}\Gamma_{\beta\alpha}^r \left(\overline{u}^k\Gamma^r d^m\right) \to -\frac{1}{4}\Gamma_{\beta\alpha}^r \left(\overline{u}\Gamma^r G^{\rho\delta}d\right).$$

This recipe enables us to insert the gluon field strength tensor G into quark matrix elements, and leads to generation of three-particle local matrix elements of the pion. We neglect the terms $\sim G^2$ in our computations.

Having finished a color summation one can calculate the traces over spinor indices and perform integrations to extract the imaginary part of the correlation function $\Pi^{\mathrm{QCD}}(p,q)$ in accordance with procedures described in Ref. [22]. Omitting the technical details we provide the final expression for the spectral density, which consists of the perturbative and nonperturbative components

$$\rho^{\mathrm{QCD}}(s) = \rho^{\mathrm{pert.}}(s) + \rho^{\mathrm{n.-pert.}}(s). \tag{16}$$

Calculations show that, the local matrix element of the pion, which contributes in the soft-meson limit to the $\mathrm{Im}\Pi^{\mathrm{QCD}}(p)$ is

$$\langle 0|\overline{d}(0)i\gamma_5 u(0)|\pi(q)\rangle = f_\pi \mu_\pi. \tag{17}$$

Here

$$\mu_\pi = \frac{m_\pi^2}{m_u + m_d} = -\frac{2\langle\overline{q}q\rangle}{f_\pi^2}, \tag{18}$$

where the second equality is the relation between m_π, f_π, the quark masses and the quark condensate $\langle \bar{q}q \rangle$, which follows from the partial conservation of axial vector current (PCAC).

The components of the spectral density are given by the formulas:

$$\rho^{\text{pert.}}(s) = \frac{f_\pi \mu_\pi}{4\pi^2 s} \sqrt{s(s - 4m_b^2)} \left(s + 2m_b m_s - 2m_b^2 \right), \tag{19}$$

and

$$\rho^{\text{n.-pert.}}(s) = \frac{f_\pi \mu_\pi}{18} \left\{ 6\langle \bar{s}s \rangle \left[-2m_b \delta(s - m_b^2) \right. \right.$$
$$\left. + sm_s \delta^{(1)}(s - m_b^2) \right] + \langle \bar{s}g\sigma Gs \rangle \left[6(m_b - m_s)\delta^{(1)}(s - m_b^2) \right.$$
$$\left. \left. - 3s(m_b - 2m_s)\delta^{(2)}(s - m_b^2) - s^2 m_s \delta^{(3)}(s - m_b^2) \right] \right\}. \tag{20}$$

In Eq. (20) $\delta^{(n)}(s - m_b^2) = (d/ds)^n \delta(s - m_b^2)$ that appear in extracting the imaginary part of the pole terms and stem from the well known formula

$$\frac{1}{s - m_b^2} = \text{P.V.} \frac{1}{s - m_b^2} + i\pi\delta(s - m_b^2).$$

The continuum subtraction is performed in a standard manner after $\rho^h(s) \to \rho^{\text{QCD}}(s)$ replacement. Then, the final sum rule to evaluate the strong coupling reads

$$g_{X_b B_s \pi} = \frac{(m_b + m_s)}{f_{B_s} f_{X_b} m_{X_b} m_{B_s}^2 m^2} \left(1 - M^2 \frac{d}{dM^2} \right) M^2$$
$$\times \int_{(m_b + m_s)^2}^{s_0} ds\, e^{(m^2 - s)/M^2} \rho^{\text{QCD}}(s). \tag{21}$$

The width of the decay $X_b \to B_s^0 \pi^+$ can be found applying the usual prescriptions and definitions for the strong coupling together with other matrix elements from Eq. (5) as well as the parameters of the X_b state. The calculations give

$$\Gamma\left(X_b \to B_s^0 \pi^+ \right) = \frac{g_{X_b B_s \pi}^2 m_{B_s}^2}{24\pi} \lambda\left(m_{X_b},\ m_{B_s}, m_\pi \right)$$
$$\times \left[1 + \frac{\lambda^2\left(m_{X_b},\ m_{B_s}, m_\pi \right)}{m_{B_s}^2} \right], \tag{22}$$

where

$$\lambda(a,\ b,\ c) = \frac{\sqrt{a^4 + b^4 + c^4 - 2\left(a^2 b^2 + a^2 c^2 + b^2 c^2 \right)}}{2a}.$$

Equations (21) and (22) are final expressions that will be used for numerical analysis of the decay channel $X_b \to B_s^0 \pi^+$.

III. NUMERICAL RESULTS AND CONCLUSIONS

The QCD sum rule for the strong coupling $g_{X_b B_s \pi}$ and decay width $\Gamma\left(X_b \to B_s^0 \pi^+ \right)$ contain various parameters

Parameters	Values
m_{B_s}	(5366.77 ± 0.24) MeV
f_{B_s}	(222 ± 12) MeV
m_b	(4.18 ± 0.03) GeV
m_s	(95 ± 5) MeV
$\langle \bar{q}q \rangle$	$(-0.24 \pm 0.01)^3$ GeV^3
$\langle \bar{s}s \rangle$	$0.8\, \langle \bar{q}q \rangle$
m_0^2	(0.8 ± 0.1) GeV^2
$\langle \bar{q}g\sigma Gq \rangle$	$m_0^2 \langle \bar{q}q \rangle$

TABLE I: Input parameters used in calculations.

that should be fixed in accordance with the standard procedures: for numerical computations we need the masses and decay constants of the X_b and B_s^0 mesons as well as values of the quark and mixed condensates. In addition to that, QCD sum rules depend on the b and s quark masses. The values of some used parameters are moved to Table I. In the calculations we also employ the QCD sum rule predictions for the mass and decay constant of X_b state obtained in our work Ref. [14]. The value of the decay constant f_{B_s} is borrowed from Ref. [24].

Calculations also require fixing of the auxiliary parameters, namely the continuum threshold s_0 and Borel parameter M^2. The standard criteria accepted in the sum rule calculations require the practical independence of the physical quantities on these auxiliary parameters. The continuum threshold is not totally arbitrary but, in principle, depends on the energy of the first excited state with the same quantum numbers and structure as the particle under consideration. In the lack of information on the mass of the first excited state in this channel, we follow the traditional prescriptions and choose s_0 in the interval $(m_{X_b} + 0.3)^2$ GeV$^2 \leq s_0 \leq (m_{X_b} + 0.5)^2$ GeV^2, i.e.

$$34.4 \text{ GeV}^2 \leq s_0 \leq 36.8 \text{ GeV}^2. \tag{23}$$

To determine the working window for the Borel parameter, we demand not only sufficient suppression of the contributions due to the higher states and continuum, but also exceeding of the perturbative contributions over the non-perturbative ones as well as convergence of the OPE series. Technically, the upper limit on M^2 is obtained by the requirement

$$\frac{\int_0^{s_0} ds \rho^{QCD}(s) e^{-s/M^2}}{\int_0^\infty ds \rho^{QCD}(s) e^{-s/M^2}} > 1/2. \tag{24}$$

The lower limit on M^2 is found by requiring that the perturbative contribution exceeds the non-perturbative one, and that the higher dimensional terms constitute less than 10% of the total contribution. These requirements lead to the working interval 6 GeV$^2 \leq M^2 \leq 8$ GeV^2 for the Borel parameter. Considering these regions, we plot the strong coupling constant $g_{X_b B_s \pi}$ as functions of M^2 and s_0 in Figs. 1 and 2. From these figures we see
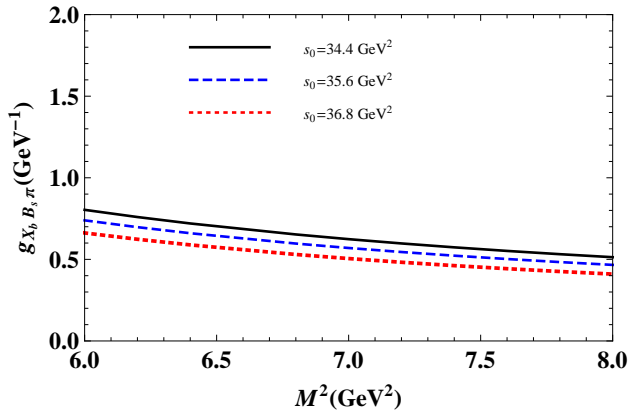
FIG. 1: The strong coupling constant $g_{X_b B_s \pi}$ vs Borel parameter M^2 at different fixed values of s_0.
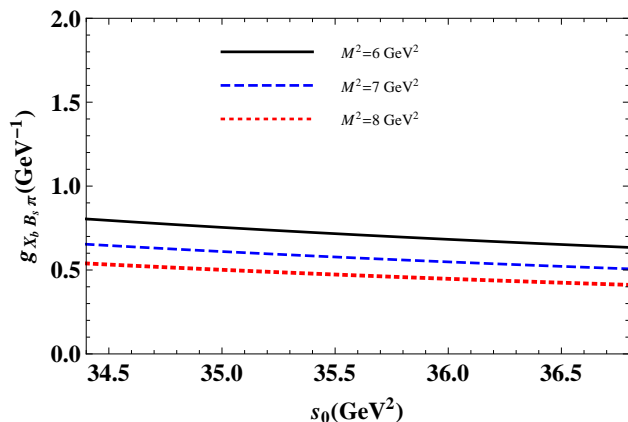

FIG. 2: The strong coupling constant $g_{X_b B_s \pi}$ as a function of the threshold parameter s_0 at different fixed values of M^2.

that the coupling $g_{X_b B_s \pi}$ demonstrates weak dependence on the Borel and threshold parameters in the selected working regions.

Extracted from the numerical calculations, value of the strong coupling $g_{X_b B_s \pi}$ is obtained as

$$g_{X_b B_s \pi} = (0.60 \pm 0.23) \text{ GeV}^{-1}. \qquad (25)$$

For the width of the decay $X_b(5568) \to B_s^0 \pi^+$ we get

$$\Gamma(X_b \to B_s^0 \pi^+) = (22.4 \pm 9.2) \text{ MeV}, \qquad (26)$$

which is in a good consistency with the experimental data of the D0 Collaboration.

In this work we have continued our studies of the newly discovered exotic state $X_b(5568)$ and computed the width of the strong decay $X_b \to B_s^0 \pi^+$ using methods of QCD sum rules on the light-cone and soft-meson approximation. To this end, first we found the strong coupling $g_{X_b B_s \pi}$ that allowed us to calculate $\Gamma(X_b \to B_s^0 \pi^+)$. Our finding is consistent with the experimental data of the D0 Collaboration.

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