Analysis of heavy spin-3/2 baryon-heavy spin-1/2 baryon-light vector meson vertices in QCD

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The heavy spin-3/2 baryon-heavy spin-1/2 baryon vertices with light vector mesons are studied within the light cone QCD sum rules method. These vertices are parametrized in terms of three coupling constants. These couplings are calculated for all possible transitions. It is shown that correlation functions for these transitions are described by only one invariant function for every Lorenz structure. The obtained relations between the correlation functions of the different transitions are structure independent while explicit expressions of invariant functions depend on the Lorenz structure.

I. INTRODUCTION

In the last decade, significant experimental progress has been achieved in heavy baryon physics. These highly excited (and often unexpected) experimental results have been announced by the *BaBar*, Belle, CDF and D0 Collaborations. The $\frac{1}{2}^+$ and $\frac{1}{2}^-$ antitriplet states, Λ_c^+, Ξ_c^+, Ξ_c^0 and $\Lambda_c^+(2593)$, $\Xi_c^+(2790)$, $\Xi_c^0(2790)$ as well as the $\frac{1}{2}^+$ and $\frac{3}{2}^+$ sextet states, Ω_c, Σ_c, Ξ_c' and Ω_c^*, Σ_c^*, Ξ_c^* have been observed [1,2]. Among the S-wave bottom baryons, only Λ_b, Σ_b, Σ_b^*, Ξ_b and Ω_b have been discovered. Moreover, the physics at LHC opens new horizons for detailed study of the observed heavy baryons, and paves the way to an invaluable opportunity for search of the new baryon states [3]. Evaluation of the strong heavy baryon-heavy baryon-meson coupling constants could be important in analysis of $\bar{p}p$ and e^+e^- experiments with pair production of heavy baryons with spin-1/2 or spin-3/2. For example, in *BaBar* and BELLE, these baryons can be produced in pairs, where one of them is off shell. Analysis of the subsequent strong decays of these baryons requires knowledge about their strong coupling constants.

Considerable progress in experiments has stimulated theoretical analysis of the heavy flavor physics. Heavy baryons with a single heavy quark can serve as an excellent "laboratory" for testing predictions of the quark models and heavy quark symmetry. After discovery of heavy baryons with a single heavy quark the next step of investigations in this direction is to study their strong, electromagnetic and weak decays which can give us useful information on the quark structure of these baryons.

The strong coupling constants of these baryons with light vector mesons are the main ingredients for their

strong decays and a more accurate determination of these constants is needed. For this aim one should consult some kind of nonperturbative methods in QCD as we deal at the hadronic scale. The method of the QCD sum rules [4] proved to be one of the most predictive among all other nonperturbative methods, and in this respect, the most advanced version seems to be the formalism implemented on the light cone. In the light cone QCD sum rules (LCSR), the operator product expansion (OPE) is carried out near the light cone, $x^2 \sim 0$, and the nonperturbative hadronic matrix elements are parametrized in terms of distribution amplitudes (DA's) of a given particle (for more about LCSR, see [5]).

In [6–8], the strong coupling constants of light pseudoscalar and vector mesons with a sextet and antitriplet of the spin-1/2 heavy baryons, as well as the heavy spin-3/2 baryon-heavy spin-1/2 baryon vertices with light pseudoscalar mesons, are calculated within the light cone version of the QCD sum rules. In the present work, we extend our previous studies to investigate the strong coupling constants among sextet of the heavy spin-3/2 baryons and the sextet and antitriplet of the heavy spin-1/2 baryons and the light vector mesons.

The plan of this paper is as follows. We first derive LCSR for the coupling constants of the transitions of the sextet spin-3/2 heavy baryons to sextet and antitriplet spin-1/2 heavy baryons and light vector mesons. In Sec. III, we present our numerical analysis of the aforementioned coupling constants and compare our predictions with the results available on this subject.

II. LIGHT CONE QCD SUM RULES FOR $B_Q^* B_Q V$ VERTICES

In this section, we calculate the strong coupling constants $B_Q^* B_Q^6 V$ and $B_Q^* B_Q^3 V$, where B_Q^* is the heavy spin-3/2 sextet, B_Q^6 stands for the heavy spin-1/2 sextet and B_Q^3 denotes the heavy spin-1/2 antitriplet baryons. The vertex describing spin-3/2 baryon transition into spin-1/2

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baryon and light vector meson can be parametrized in the following way [9]:

$$\langle B_Q(p_2)V(q)|B_Q^*(p_1)\rangle$$

$$= \bar{u}_{B_Q}(p_2)\Big\{g_1(q_\alpha \not\varepsilon - \varepsilon_\alpha \not q)\gamma_5 + g_2[(P \cdot \varepsilon)q_\alpha$$

$$- (P \cdot q)\varepsilon_\alpha]\gamma_5 + g_3[(q \cdot \varepsilon)q_\alpha - q^2\varepsilon_\alpha]\gamma_5\Big\}u_{B_{Q\alpha}^*}(p_1),$$

$$(1)$$

where $u_{B_Q}(p_2)$ is the Dirac spinor of either the sextet baryon B_Q^6 or the antitriplet baryon B_Q^3, while $u_{B_{Q\alpha}^*}(p_1)$ is the Rarita-Schwinger spinor of the spin-3/2 sextet baryon B_Q^*, ε_μ is the polarization 4-vector of the light vector meson and $2P = p_1 + p_2$, $q = p_1 - p_2$. Furthermore, the conditions $q^2 = m_V^2$ and $(q\varepsilon) = 0$ are imposed for the on shell vector meson. Here we would like to make the following remark. Many of considered transitions in this work are kinematically forbidden. In other words, one

of the particles should be off the mass shell and therefore "coupling constants" have q^2 dependence. We calculate these form factors at $q^2 = m_V^2$ and assume that in going q^2 from this point to q_{max}^2 (in our case $q_{max}^2 = (m_1 - m_2)^2$ indeed is small), the coupling constants do not change considerably (for a detailed discussion see [10]).

In order to calculate the strong coupling constants g_k, $k = 1, 2, 3$, in the framework of the LCSR, we start by considering the correlation function:

$$\Pi_\mu^{B_Q^* \to B_Q V} = i \int d^4x e^{ip_2 x}\Big\langle V(q)\Big|T\{\eta^{B_Q}(x)\bar{\eta}_\mu^{B_Q^*}(0)\}\Big|0\Big\rangle,$$

$$(2)$$

where η_μ and η are interpolating currents of the B_Q^* or $B_Q^{6,3}$ baryons, respectively, with T being the time ordering operator.

The general form of the interpolating current of the spin-1/2 heavy sextet B_b^6 and antitriplet B_b^3 baryons can be written in the following form [11]:

$$\eta_Q^{(6)}(q_1, q_2, Q) = -\sqrt{\frac{1}{2}}\epsilon^{abc}\Big[\Big(q_1^{aT}CQ^b\Big)\gamma_5 q_2^c - \Big(Q^{aT}Cq_2^b\Big)\gamma_5 q_1^c + \beta\Big(q_1^{aT}C\gamma_5 Q^b\Big)q_2^c - \beta\Big(Q^{aT}C\gamma_5 q_2^b\Big)q_1^c\Big],$$

$$\eta_Q^{(3)}\Big(q_1, q_2, Q\Big) = \sqrt{\frac{1}{6}}\epsilon^{abc}\Big[2\Big(q_1^{aT}Cq_2^b\Big)\gamma_5 Q^c + \Big(q_1^{aT}CQ^b\Big)\gamma_5 q_2^c + \Big(Q^{aT}Cq_2^b\Big)\gamma_5 q_1^c + 2\beta(q_1^{aT}C\gamma_5 q_2^b)Q^c$$

$$+ \beta\Big(q_1^{aT}C\gamma_5 Q^b\Big)q_2^c + \beta\Big(Q^{aT}C\gamma_5 q_2^b\Big)q_1^c\Big],$$

$$(3)$$

where a, b, c are the color indices, C is the charge conjugation operator, β is an arbitrary parameter and $\beta = -1$ corresponds to the choice for the Ioffe current [12]. The interpolating current for the spin-3/2 sextet baryons can be written as [11]

$$\eta_\mu = A\epsilon^{abc}\Big[\Big(q_1^{aT}C\gamma_\mu q_2^b\Big)Q^c + \Big(q_2^{aT}C\gamma_\mu Q^b\Big)q_1^c$$

$$+ \Big(Q^{aT}C\gamma_\mu q_1^b\Big)q_2^c\Big].$$

$$(4)$$

The light quark contents of both sextet and antitriplet heavy spin-1/2 baryons are shown in Table I. The values

of normalization constant A and the quark flavors q_1, q_2 and Q for each member of the heavy spin-3/2 baryon sextet are also listed in Table II.

Using the quark-hadron duality and inserting a complete set of hadronic states with the same quantum numbers as the interpolating currents η_μ and $\eta^{6,3}$ into a correlation function, one can obtain the representation of $\Pi_\mu(p, q)$ in terms of hadrons. Isolating the ground state contributions coming from the heavy baryons in the corresponding channels we get

TABLE I. The light quark content q_1 and q_2 for the sextet and antitriplet baryons with spin-1/2.

	$\Sigma_{b(c)}^{+(++)}$	$\Sigma_{b(c)}^{0(+)}$	$\Sigma_{b(c)}^{-(0)}$	$\Xi_{b(c)}^{-(0)'}$	$\Xi_{b(c)}^{0(+)'}$	$\Omega_{b(c)}^{-(0)}$	$\Lambda_{b(c)}^{0(+)}$	$\Xi_{b(c)}^{-(0)}$	$\Xi_{b(c)}^{0(+)}$
q_1	u	u	d	d	u	s	u	d	u
q_2	u	d	d	s	s	s	d	s	s

TABLE II. The light quark content q_1, q_2 and normalization constant A for the sextet baryons with spin-3/2.

	$\Sigma_{b(c)}^{*+(++)}$	$\Sigma_{b(c)}^{*0(+)}$	$\Sigma_{b(c)}^{*-(0)}$	$\Xi_{b(c)}^{*0(+)}$	$\Xi_{b(c)}^{*-(0)}$	$\Omega_{b(c)}^{*-(0)}$
q_1	u	u	d	u	d	s
q_2	u	d	d	s	s	s
A	$\sqrt{1/3}$	$\sqrt{2/3}$	$\sqrt{1/3}$	$\sqrt{2/3}$	$\sqrt{2/3}$	$\sqrt{1/3}$

$$\Pi_\mu(p, q) = \frac{\langle 0|\eta|B_Q(p_2)\rangle\langle B_Q(p_2)V(q)|B_Q^*(p_1)\rangle}{p_2^2 - m_2^2} \frac{\langle B_Q^*(p_1)|\bar{\eta}_\mu|0\rangle}{p_1^2 - m_1^2} + \cdots$$

$$= -\frac{\lambda_{B_Q}\lambda_{B_Q^*}(\not{p}_2 + m_2)}{(p_2^2 - m_2^2)(p_1^2 - m_1^2)}\Big\{g_1(q_\alpha\not{\varepsilon} - \varepsilon_\alpha\not{q})\gamma_5 + g_2[(P \cdot \varepsilon)q_\alpha - (P \cdot q)\varepsilon_\alpha]\gamma_5 + g_3[(q \cdot \varepsilon)q_\alpha - q^2\varepsilon_\alpha]\gamma_5\Big\}$$

$$\times (\not{p}_1 + m_1)\Big(g_{\alpha\mu} - \frac{\gamma_\alpha\gamma_\mu}{3m_1} - \frac{2p_{1\alpha}p_{1\mu}}{3m_1^2} + \frac{p_{1\alpha}\gamma_\mu - p_{1\mu}\gamma_\alpha}{3m_1}\Big), \tag{5}$$

where $m_1 = m_{B^*}$, $m_2 = m_{B_Q}$, and $\cdots$ represent the contributions of the higher states and continuum. In derivation of Eq. (5) we have used the following definitions:

$$\langle 0|\eta|B_Q(p)\rangle = \lambda_{B_Q}u(p, s),$$
$$\langle 0|\eta_\mu|B_Q^*(p)\rangle = \lambda_{B_Q^*}u_\mu(p, s), \tag{6}$$

and summation over the spins has been performed using the relations

$$\sum_s u(p_2, s)\bar{u}(p_2, s) = \not{p}_2 + m_2, \sum_s u_\alpha(p_1, s)\bar{u}_\beta(p_1, s)$$

$$= -(\not{p}_1 + m_1)\Big(g_{\alpha\beta} - \frac{\gamma_\alpha\gamma_\beta}{3m_1}$$

$$- \frac{2p_{1\alpha}p_{1\beta}}{3m_1^2} + \frac{p_{1\alpha}\gamma_\beta - p_{1\beta}\gamma_\alpha}{3m_1}\Big). \tag{7}$$

Here we would like to make following remark. The interpolating current η_μ couples not only to the $J^P = \frac{3}{2}^+$ states, but also to the $J^P = \frac{1}{2}^-$ states. The corresponding matrix element of the current η_μ between vacuum and $J^P = \frac{1}{2}^-$ states can be parametrized as follows:

$$\langle 0|\eta_\mu|\tilde{B}_Q(p)\rangle = \tilde{\lambda}\Big(\gamma_\mu - 4\frac{p_\mu}{\tilde{m}}\Big)u(p, s), \tag{8}$$

where a tilde means a $J^P = \frac{1}{2}^-$ state, and $\tilde{\lambda}$ and $\tilde{m}$ represent its residue and mass, respectively. Using Eqs. (5) and (8), we see that the structures proportional to γ_μ at the right end and to $p_{1\mu} = p_{2\mu} + q_\mu$ receive contributions not only from $J^P = \frac{3}{2}^+$ states but also from the $J^P = \frac{1}{2}^-$ states which should be removed.

Another problem is that not all Lorentz structures are independent. We can remove both problems by ordering the Dirac matrices in a specific way which guarantees the independence of all the Lorentz structures as well as the absence of the $J^P = \frac{1}{2}^-$ contributions. In the present work, we choose the ordering of the Dirac matrices in the form $\gamma_\mu\not{p}\not{q}\not{\varepsilon}\gamma_5$. Choosing this ordering and using Eq. (4), for the phenomenological part of the correlation function, we get

$$\Pi_\mu = \frac{\lambda_{B_Q}\lambda_{B_Q^*}}{[m_1 - (p + q)^2]} \frac{1}{(m_2^2 - p^2)}\Big[g_1(m_1 + m_2)\not{p}\not{\varepsilon}\gamma_5q_\mu$$

$$- g_2\not{p}\not{q}\gamma_5(p \cdot \varepsilon)q_\mu + g_3q^2\not{p}\not{q}\gamma_5\varepsilon_\mu$$

$$+ \text{other structures}\Big], \tag{9}$$

where we have set $p_2 = p$. In order to calculate the coupling constants g_1, g_2 and g_3 from the QCD sum rules, we should know the expression for the correlation function from the QCD side. But first we try to find relations between various invariant functions which would simplify the calculations of the constants g_1, g_2 and g_3 considerably for different channels. For this aim we will follow the approach given in [13–17] where main "construction blocks" have been derived (see also [6–8]). These relations are independent from the choice of the explicit Lorenz structure and automatically take into account violation of the flavor unitary symmetry.

We will also show that all the transitions of spin-3/2 sextet B_Q^* into the sextet and antitriplet of the spin-1/2 B_Q are described in terms of only one invariant function for each Lorenz structure. First, we consider sextet-sextet transitions and as an example we start with the $\Sigma_b^{*0} \to \Sigma_b^0\rho^0$ transition. The invariant correlation function for the $\Sigma^{*0} \to \Sigma_b^0\rho^0$ transition can be written in the general form as

$$\Pi^{\Sigma_b^{*0}\to\Sigma_b^0\rho^0} = g_{\rho^0uu}\Pi_1(u, d, b) + g_{\rho^0dd}\Pi_1'(u, d, b)$$

$$+ g_{\rho^0bb}\Pi_2(u, d, b). \tag{10}$$

The interpolating currents for Σ_b^{*0} and Σ_b^0 are symmetric with respect to the exchange of the light quarks, then obviously $\Pi_1'(u, d, b) = \Pi_1(d, u, b)$. Moreover, using Eq. (4) one can easily obtain that $-\Pi_2(u, d, b) = \Pi_1(b, u, d) + \Pi_1(b, d, u)$. Couplings of quarks to the ρ^0 meson are obtained from the quark current

$$J_\mu^{\rho^0} = \sum_{u,d,s} g_{qq\rho}\bar{q}\gamma_\mu q, \tag{11}$$

and for the ρ^0 meson $g_{\rho^0uu} = -g_{\rho^0dd} = 1/\sqrt{2}$ and similarly for ω and ϕ mesons one gets $g_{\omega uu} = g_{\omega dd} = 1/\sqrt{2}$, $g_{\phi ss} = 1$, all the other couplings to light mesons being zero.

The function $\Pi_1(u, d, b)$ describes emission of the ρ^0 meson from u, d and b quarks, respectively, and is formally defined as

$$\Pi_1(u, d, b) = \langle \bar{u}u | \Sigma_b^{*0} \bar{\Sigma}_b^{*0} | 0 \rangle. \tag{12}$$

Using Eqs. (10)–(12) we get

$$\Pi^{\Sigma_b^{*0} \to \Sigma_b^0 \rho^0} = \frac{1}{\sqrt{2}} \left[\Pi_1(u, d, b) - \Pi_1(d, u, b) \right]. \tag{13}$$

In the isospin symmetry limit, the invariant functions for the $\Sigma_b^{*0} \to \Sigma_b \rho^0$ and $\Sigma^{*0} \to \Lambda_b \omega^0$ transitions vanish, as is expected. The relations among other invariant functions involving neutral vector mesons ρ, ω and ϕ can be obtained in a similar way, and we put these relations into the Appendix.

The relations involving charged ρ mesons require some care. Indeed, in the $\Sigma_b^{*0} \to \Sigma_b^0 \rho^0$ transition $u(d)$ quarks from baryons Σ_b^{*0} and Σ_b^0 form the $\bar{u}u(\bar{d}d)$ state, while d (u) and b quarks are spectators. In the case of charged ρ^+ meson, d quark from baryons Σ_b^0 and u quark from baryon Σ_b^{*0} form the $(\bar{u}d)$ state while the remaining $d(u)$ b quarks again are spectators. Therefore it is quite natural to expect that these matrix elements should be proportional. Indeed, explicit calculations confirm this expectation and we find that

$$\Pi^{\Sigma_b^{*+} \to \Sigma_b^0 \rho^+} = \langle \bar{d} | \Sigma_b^0 \bar{\Sigma}_b^{*+} | 0 \rangle = -\sqrt{2} \langle \bar{d}d | \Sigma_b^0 \bar{\Sigma}_b^{*0} | 0 \rangle$$
$$= -\sqrt{2} \Pi_1(d, u, b). \tag{14}$$

Replacing $u \leftrightarrow d$ in Eq. (14), we get

$$\Pi^{\Sigma_b^{*-} \to \Sigma_b^0 \rho^-} = -\sqrt{2} \Pi_1(u, d, b). \tag{15}$$

It should be noted that the relations between invariant functions involving ρ and ω mesons can also be obtained from the isotopic symmetry argument. All other relations among invariant functions involving charged ρ, K^* and ϕ mesons are obtained in a similar way with the proper change of quark symbols and are presented in the Appendix. To describe the sextet-sextet transitions, we need to calculate the invariant function Π_1. For this aim, the correlation function which describes the transition $\Sigma^{*0} \to \Sigma^0 \rho^0$ would serve as a good candidate.

Up to now we have discussed the sextet-sextet transitions and found that all these transitions involving light vector mesons are described by a single universal function. We now proceed discussing the sextet-antitriplet transitions. Our goal here is to show that these transitions, similar to the sextet-sextet, are also described with the same invariant function. For this aim let us consider the $\Sigma^{*0} \to \Lambda_b \rho^0$ transition.

Similar to Eq. (10), this transition can be written as

$$\Pi^{\Sigma^{*0} \to \Lambda \rho^0} = g_{\rho^0 uu} \tilde{\Pi}_1(u, d, b) + g_{\rho^0 dd} \tilde{\Pi}_1'(u, d, b)$$
$$+ g_{\rho^0 bb} \tilde{\Pi}_2(u, d, b), \tag{16}$$

where a tilde is used to note the difference of the invariant function responsible for the sextet-antitriplet transition from the sextet-sextet transition. In order to express the $\tilde{\Pi}$ in terms of Π, let us first express the interpolating current of Λ_b in terms of sextet current. Performing similar calculations as is done in [16], the following relation between the two currents can easily be obtained:

$$\eta^{(6)}(q_1 \leftrightarrow Q) - \eta^{(6)}(q_2 \leftrightarrow Q) = \sqrt{3} \eta^{(3)}(q_1, q_2, Q),$$
$$\eta^{(6)}(q_1 \leftrightarrow Q) + \eta^{(6)}(q_2 \leftrightarrow Q) = -\eta^{(6)}(q_1, q_2, Q). \tag{17}$$

Using these relations and Eq. (10), we construct the following auxiliary quantities:

$$\Pi^{\Sigma_b^{*0} \to \Sigma_b(u \leftrightarrow b) \rho^0} = g_{\rho bb} \Pi_1(b, d, u) + g_{\rho dd} \Pi_1'(b, d, u)$$
$$+ g_{\rho uu} \Pi_2(b, d, u), \tag{18}$$

$$\Pi^{\Sigma_b^{*0} \to \Sigma_b(d \leftrightarrow b) \rho^0} = g_{\rho uu} \Pi_1(u, b, d) + g_{\rho bb} \Pi_1'(u, b, d)$$
$$+ g_{\rho dd} \Pi_2(u, b, d). \tag{19}$$

From these expressions, we immediately obtain that

$$\sqrt{3} \Pi^{\Sigma_b^{*0} \to \Lambda_b \rho^0} = g_{\rho uu} \left[\Pi_2(b, d, u) - \Pi_1(u, b, d) \right]$$
$$+ g_{\rho dd} \left[\Pi_1'(b, d, u) - \Pi_2(u, b, d) \right]$$
$$+ g_{\rho bb} \left[\Pi_1(b, d, u) - \Pi_1'(u, b, d) \right], \tag{20}$$

$$-\Pi^{\Sigma_b^{*0} \to \Sigma_b^0 \rho^0} = g_{\rho uu} \left[\Pi_2(b, d, u) + \Pi_1(u, b, d) \right]$$
$$+ g_{\rho dd} \left[\Pi_1'(b, d, u) + \Pi_2(u, b, d) \right]$$
$$+ g_{\rho bb} \left[\Pi_1(b, d, u) + \Pi_1'(u, b, d) \right]$$
$$= -g_{\rho uu} \Pi_1(u, d, b) - g_{\rho dd} \Pi_1'(u, d, b)$$
$$- g_{\rho bb} \Pi_2(u, d, b), \tag{21}$$

where in obtaining the last line, we have used Eq. (10). From this equation, we immediately get

$$-\Pi_2(b, d, u) = \Pi_1(u, b, d) + \Pi_1(u, d, b), \tag{22}$$

$$-\Pi_2(u, b, d) = \Pi_1'(b, d, u) + \Pi_1'(u, d, b), \tag{23}$$

$$-\Pi_2(u, d, b) = \Pi_1(b, d, u) + \Pi_1'(u, b, d). \tag{24}$$

With the replacement $b \leftrightarrow u$ Eq. (22) goes to Eq. (24) and with the replacement $b \leftrightarrow d$ Eq. (23) goes to Eq. (24), as the result of which we get

$$\Pi_1'(u, b, d) = \Pi_1(b, d, u), \qquad \text{and,} \tag{25}$$

$$\Pi_1'(d, b, u) = \Pi_1(b, d, u). \tag{26}$$

Using these relations and Eq. (20), we get the following relation for the invariant function responsible for the $\Sigma_b^{*0} \to \Lambda_b \rho^0$ transition:

$$\sqrt{3}\Pi^{\Sigma_b^{*0} \to \Lambda_b \rho^0} = -g_{\rho uu}\left[2\Pi_1(u,b,d) + \Pi_1(u,d,b)\right]$$
$$+ g_{\rho dd}\left[2\Pi_1(d,b,u) + \Pi_2(d,u,b)\right]$$
$$+ g_{\rho bb}\left[\Pi_1(b,d,u) - \Pi_1(b,u,d)\right]. \quad (27)$$

Comparing these results with Eq. (16), we finally get

$$\tilde{\Pi}_1(u,d,b) = -\frac{1}{\sqrt{3}}\left[2\Pi_1(u,b,d) + \Pi_1(u,d,b)\right],$$
$$\tilde{\Pi}'_1(u,d,b) = \frac{1}{\sqrt{3}}\left[2\Pi_1(d,b,u) + \Pi_1(d,u,b)\right], \quad (28)$$
$$\tilde{\Pi}_2(u,d,b) = \frac{1}{\sqrt{3}}\left[\Pi_1(b,d,u) - \Pi_1(b,u,d)\right].$$

Relations among invariant functions describing sextet-antitriplet transitions involving light vector mesons are presented in the Appendix.

For obtaining sum rules for the coupling constants the expressions of the correlation functions from the QCD side are needed. The corresponding correlation functions can be evaluated in the deep Euclidean region, $-p_1^2 \to \infty$, $-p_2^2 \to \infty$, as has already been mentioned, using the OPE. In the Light Cone version of the QCD sum rules formalism, the OPE is performed with respect to twists of the corresponding nonlocal operators. In this expansion the DA's of the vector mesons appear as the main nonperturbative parameters. Up to twist-4 accuracy, matrix elements $\langle V(q)|\bar{q}(x)\Gamma q(0)|0\rangle$ and $\langle V(q)|\bar{q}(x)G_{\mu\nu}q(0)|0\rangle$ are determined in terms of the DA's of the vector mesons, where Γ represents the Dirac matrices relevant to the case under consideration, and $G_{\mu\nu}$ is the gluon field strength tensor. The definitions of these DA's for vector mesons are presented in [18,19]. Having the expressions of the heavy and light quark propagators (see [20,21]) and the DA's for the light vector mesons we can straightforwardly calculate the correlation functions from the QCD side. Equating both representations of the correlation function and separating coefficients of Lorentz structures $\not{p}\not{\varepsilon}\gamma_5 q_\mu$, $\not{p}\not{q}\gamma_5(p\cdot\varepsilon)q_\mu$ and $\not{p}\not{q}\gamma_5\varepsilon_\mu$, and applying Borel transformation to the variables p^2 and $(p+q)^2$ on both sides of the correlation functions, which suppresses the contributions of the higher states and continuum, we obtain sum rules for the coupling constants g_1, g_2 and g_3:

$$g_k = \kappa_k \Pi_1^{(k)}(M^2), \quad (29)$$

where

$$\kappa_1 = \frac{\kappa_2}{m_1 + m_2},$$
$$\kappa_2 = \frac{1}{\lambda_{B_Q}\lambda_{B_Q^*}}e^{\left(\frac{m_1^2}{M_1^2}+\frac{m_2^2}{M_2^2}+\frac{m_V^2}{M_1^2+M_2^2}\right)}, \quad (30)$$
$$\kappa_3 = \frac{\kappa_2}{m_V^2}.$$

As has already been noted, relations between invariant functions are independent of the Lorenz structures, but their explicit expressions are structure dependent. So we have introduced extra upper index k in brackets for each coupling, and $k = 1$, 2 and 3 corresponds the choice of the Lorenz structures $\not{p}\not{\varepsilon}\gamma_5 q_\mu$, $\not{p}\not{q}\gamma_5(p\cdot\varepsilon)q_\mu$ and $\not{p}\not{q}\gamma_5\varepsilon_\mu$, respectively. In this equation M_1^2 and M_2^2 are Borel parameters in the initial and final baryon channels. Since the masses of the initial and final baryons are close to each other we put $M_1^2 = M_2^2 = M^2$. Residues $\lambda_{B_Q^*}$ and λ_{B_Q} of the heavy baryons of spin-3/2 and -1/2 have been calculated in [22]. As the explicit formulas for $\Pi_1^{(k)}$, $k = 1$, 2, 3 are lengthy and not very instructive we do not present it in the body of our article.

III. NUMERICAL ANALYSIS

In this section, we present our numerical results on the strong coupling constants of the light vector mesons with the sextet and antitriplet of heavy baryons. The main input parameters in LCSR for the coupling constants are the DA's of the light vector mesons. These DA's and its parameters are taken from [18,19]. The LCSR's also contain the following auxiliary parameters: Borel parameter M^2, threshold of the continuum s_0 and the parameter β of the interpolating currents of the spin-1/2 baryons. Obviously any physical quantity should be independent of these auxiliary parameters. Therefore, we should find "working regions" of these parameters where coupling constants of the $B_Q^* B_Q V$ transitions are practically independent of them. We proceed along the same scheme as is presented in [13–17]. In order to find the "working region" of M^2, we require that the continuum and higher state contributions should be less then half of the dispersion integral while the contribution of the higher terms proportional to $1/M^2$ be less then 25% of the total result. These two requirements give the "working region" of M^2 in the range $15 \text{ GeV}^2 \le M^2 \le 30 \text{ GeV}^2$ for baryons with the single b quark and $4 \text{ GeV}^2 \le M^2 \le 8 \text{ GeV}^2$ for the charmed baryons, respectively. The continuum threshold is not totally arbitrary but is correlated to the energy of the first excited states with the same quantum numbers as the interpolating currents. This parameter is chosen in the range $(m_{B_Q^*} + 0.5 \text{ GeV})^2 \le s_0 \le (m_{B_Q^*} + 0.7 \text{ GeV})^2$. Our results show weak dependence on this parameter in this working region.
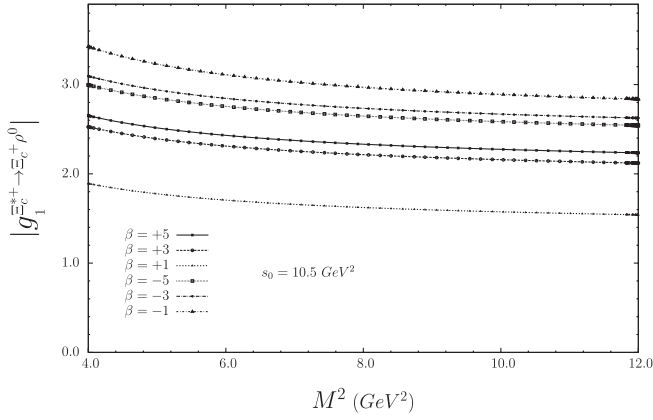
FIG. 1. The dependence of the strong coupling constant g_1 for the $\Xi_c^{*+} \rightarrow \Xi_c'^+ \rho^0$ transition on the Borel mass parameter M^2 at several different fixed values of β, and at $s_0 = 10.5$ GeV^2.

FIG. 4. The dependence of the strong coupling constant g_1 for the $\Xi_c^{*+} \rightarrow \Xi_c'^+ \rho^0$ transition on $\cos\theta$ at several different fixed values of s_0, and at $M^2 = 8.0$ GeV^2.

As an example let us consider the transition $\Xi_c^{*+} \rightarrow \Xi_c'^+ \rho^0$ and show in what way the coupling constants g_1, g_2 and g_3 are determined. In Figs. 1–3 we depict the dependence of the coupling constants g_1, g_2 and g_3 on M^2 at $s_0 = 10.5$ GeV^2 and several different fixed values of β.

It is seen that the coupling constants depend weakly on M^2 in the "working region". Now, we proceed to calculate the working region of the general parameter β entering the interpolating currents of the spin-$1/2$ particles. This parameter is also not physical quantity, hence we should

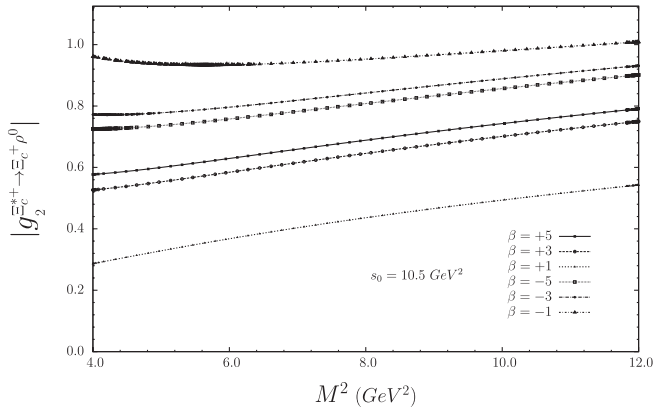FIG. 2. The same as Fig. 1, but for the strong coupling constant g_2.

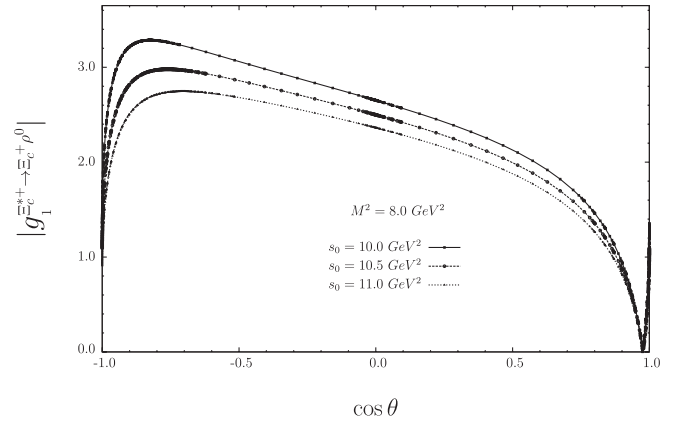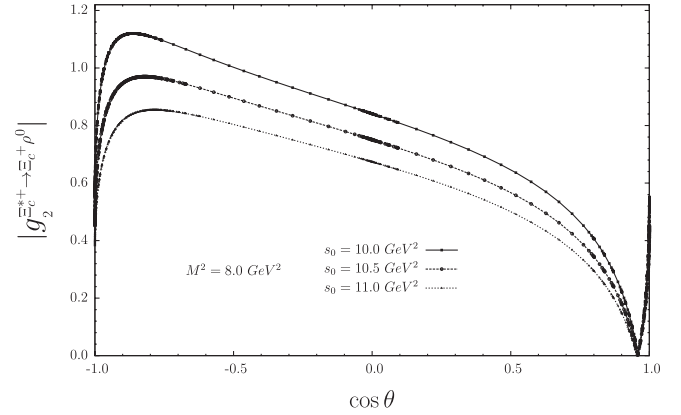FIG. 5. The same as Fig. 4, but for the strong coupling constant g_2.
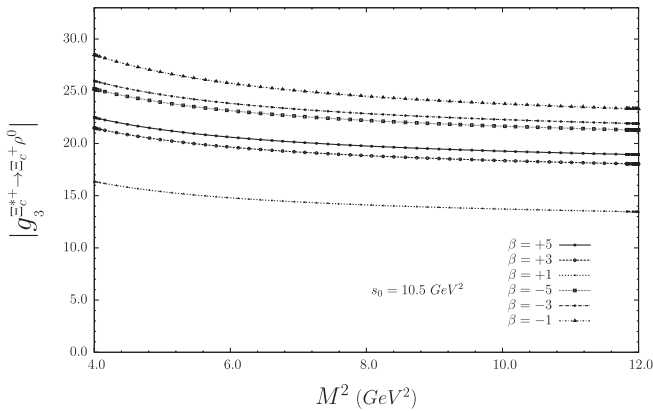
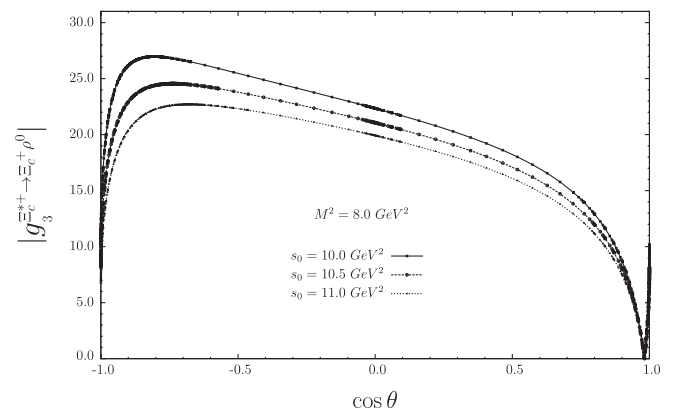FIG. 3. The same as Fig. 1, but for the strong coupling constant g_3.

FIG. 6. The same as Fig. 4, but for the strong coupling constant g_3.

TABLE III. The absolute values of the coupling constants $g_{1,2,3}$ for transitions of b-baryons. The couplings, g_1, g_2 and g_3 are in GeV^{-1}, GeV^{-1} and GeV^{-1}, respectively.

transition	g_1	g_1^{Ioffe}	g_2	g_2^{Ioffe}	g_3	g_3^{Ioffe}	NRQM
$\Sigma_b^{*0} \to \Sigma_b^- \rho^+$	4.2 ± 1.0	4.6 ± 1.2	0.7 ± 0.2	0.7 ± 0.2	84 ± 20	90 ± 23	$(2/3)c$
$\Sigma_b^{*+} \to \Sigma_b^+ \omega$	3.8 ± 1.1	4.2 ± 1.1	0.6 ± 0.2	0.7 ± 0.2	70 ± 20	74 ± 21	$(2/3)c$
$\Sigma_b^{*-} \to \Lambda_b \rho^-$	8.0 ± 2.1	8.4 ± 1.4	0.9 ± 0.3	0.8 ± 0.2	148 ± 38	154 ± 26	$(2/\sqrt{3})c$
$\Sigma_b^{*+} \to \Xi_b^{\prime 0} K^{*+}$	4.6 ± 1.4	5.2 ± 1.3	0.9 ± 0.3	1.2 ± 0.4	64 ± 18	70 ± 18	$(2/3)c$
$\Sigma_b^{*0} \to \Xi_b^0 K^{*0}$	6.4 ± 1.7	6.8 ± 1.2	0.9 ± 0.2	1.1 ± 0.2	87 ± 20	88 ± 15	$(\sqrt{2/3})c$
$\Xi_b^{*0} \to \Xi_b^{\prime 0} \rho^0$	2.4 ± 0.7	2.7 ± 0.7	0.4 ± 0.1	0.4 ± 0.1	48 ± 10	52 ± 14	$(1/3)c$
$\Xi_b^{*0} \to \Xi_b^{\prime 0} \omega$	2.4 ± 0.5	2.4 ± 0.6	0.4 ± 0.1	0.3 ± 0.1	40 ± 10	42 ± 12	$(1/3)c$
$\Xi_b^{*0} \to \Xi_b^{\prime 0} \phi$	3.2 ± 1.0	3.6 ± 1.0	0.3 ± 0.1	0.4 ± 0.1	35 ± 11	39 ± 12	$(\sqrt{2}/3)c$
$\Xi_b^{*0} \to \Sigma_b^+ K^{*-}$	4.7 ± 1.4	4.9 ± 1.4	1.0 ± 0.3	1.2 ± 0.4	64 ± 16	66 ± 16	$(2/3)c$
$\Xi_b^{*0} \to \Omega_b^- K^{*+}$	5.3 ± 1.6	5.9 ± 1.7	1.0 ± 0.2	1.3 ± 0.3	74 ± 21	83 ± 20	$(2/3)c$
$\Xi_b^{*0} \to \Xi_b^0 \rho^0$	4.3 ± 1.1	4.4 ± 0.8	0.4 ± 0.1	0.4 ± 0.1	82 ± 20	84 ± 16	$(1/\sqrt{3})c$
$\Xi_b^{*0} \to \Xi_b^0 \omega$	4.3 ± 1.1	4.4 ± 0.8	0.4 ± 0.1	0.4 ± 0.1	82 ± 20	84 ± 16	$(1/\sqrt{3})c$
$\Xi_b^{*0} \to \Xi_b^0 \phi$	5.9 ± 1.6	6.4 ± 1.1	0.7 ± 0.2	0.7 ± 0.2	65 ± 17	69 ± 12	$(\sqrt{2/3})c$
$\Xi_b^{*0} \to \Lambda_b \bar{K}^{*0}$	6.3 ± 1.7	6.6 ± 1.0	0.5 ± 0.1	0.6 ± 0.2	86 ± 20	88 ± 15	$(\sqrt{2/3})c$
$\Omega_b^{*-} \to \Omega_b^- \phi$	7.2 ± 1.8	8.2 ± 2.2	1.0 ± 0.2	1.2 ± 0.3	85 ± 15	90 ± 25	$(2\sqrt{2}/3)c$
$\Omega_b^{*-} \to \Xi_b^{\prime 0} K^{*-}$	5.3 ± 1.5	5.8 ± 1.6	1.0 ± 0.3	1.2 ± 0.4	74 ± 18	77 ± 15	$(2/3)c$
$\Omega_b^{*-} \to \Xi_b^- \bar{K}^{*0}$	9.7 ± 2.5	10.0 ± 2.0	1.3 ± 0.3	1.5 ± 0.3	136 ± 31	133 ± 25	$(2/\sqrt{3})c$

TABLE IV. The absolute values of the coupling constants $g_{1,2,3}$ for transitions of charmed baryons. The couplings, g_1, g_2 and g_3 are in GeV^{-1}, GeV^{-2} and GeV^{-2}, respectively.

'transition	g_1	g_1^{Ioffe}	g_2	g_2^{Ioffe}	g_3	g_3^{Ioffe}	NRQM
$\Sigma_c^{*+} \to \Sigma_c^0 \rho^+$	5.0 ± 1.0	5.7 ± 0.5	1.0 ± 0.2	1.2 ± 0.2	40 ± 7	45 ± 5	$(2/3)c$
$\Sigma_c^{*++} \to \Sigma_c^{++} \omega$	4.4 ± 0.8	5.0 ± 0.5	0.9 ± 0.2	1.1 ± 0.2	35 ± 6	39 ± 6	$(2/3)c$
$\Sigma_c^{*0} \to \Lambda_c^+ \rho^-$	9.6 ± 1.8	10.6 ± 1.8	1.5 ± 0.6	1.9 ± 0.3	75 ± 13	79 ± 10	$(2/\sqrt{3})c$
$\Sigma_c^{*++} \to \Xi_c^{\prime +} K^{*+}$	5.4 ± 1.0	6.0 ± 0.5	1.5 ± 0.5	2.3 ± 0.4	29 ± 4	32 ± 3	$(2/3)c$
$\Sigma_c^{*+} \to \Xi_c^+ K^{*0}$	7.5 ± 1.4	8.2 ± 1.3	2.2 ± 0.8	2.6 ± 0.4	41 ± 6	42 ± 4	$(\sqrt{2/3})c$
$\Xi_c^{*+} \to \Xi_c^{\prime +} \rho^0$	2.6 ± 0.5	3.0 ± 0.3	0.9 ± 0.2	1.0 ± 0.2	22 ± 4	25 ± 4	$(1/3)c$
$\Xi_c^{*+} \to \Xi_c^{\prime +} \omega$	2.3 ± 0.4	2.7 ± 0.4	0.8 ± 0.2	0.9 ± 0.2	19 ± 3	21 ± 2	$(1/3)c$
$\Xi_c^{*+} \to \Xi_c^{\prime +} \phi$	3.6 ± 0.7	4.1 ± 0.6	0.9 ± 0.2	1.1 ± 0.2	17 ± 3	20 ± 2	$(\sqrt{2}/3)c$
$\Xi_c^{*+} \to \Sigma_c^{++} K^{*-}$	5.5 ± 1.0	6.1 ± 0.8	1.8 ± 0.4	2.3 ± 0.4	29 ± 4	31 ± 5	$(2/3)c$
$\Xi_c^{*+} \to \Omega_c^0 K^{*+}$	5.6 ± 1.0	6.5 ± 1.0	2.0 ± 0.5	2.4 ± 0.5	32 ± 5	35 ± 5	$(2/3)c$
$\Xi_c^{*+} \to \Xi_c^+ \rho^0$	5.0 ± 0.9	5.4 ± 0.3	0.9 ± 0.2	1.0 ± 0.2	41 ± 7	44 ± 7	$(1/\sqrt{3})c$
$\Xi_c^{*+} \to \Xi_c^+ \omega$	5.0 ± 0.9	4.8 ± 0.6	0.9 ± 0.2	1.0 ± 0.2	41 ± 7	44 ± 7	$(1/\sqrt{3})c$
$\Xi_c^{*+} \to \Xi_c^+ \phi$	6.8 ± 1.3	7.6 ± 0.7	1.3 ± 0.3	1.6 ± 0.3	32 ± 6	34 ± 5	$(\sqrt{2/3})c$
$\Xi_c^{*+} \to \Lambda_c^+ \bar{K}^{*0}$	7.4 ± 1.3	8.0 ± 0.8	1.7 ± 0.5	2.0 ± 0.4	40 ± 5	41 ± 6	$(\sqrt{2/3})c$
$\Omega_c^{*0} \to \Omega_c^0 \phi$	7.8 ± 1.0	8.7 ± 1.0	2.1 ± 0.3	2.4 ± 0.3	37 ± 7	42 ± 4	$(2\sqrt{2}/3)c$
$\Omega_c^{*0} \to \Xi_c^{\prime +} K^{*-}$	5.8 ± 1.0	6.4 ± 0.8	2.1 ± 0.4	2.4 ± 0.3	32 ± 4	34 ± 3	$(2/3)c$
$\Omega_c^{*0} \to \Xi_c^0 \bar{K}^{*0}$	11.0 ± 2.0	12.0 ± 2.0	3.5 ± 0.6	4.4 ± 0.6	63 ± 9	65 ± 7	$(2/\sqrt{3})c$

look for an optimal working region at which the dependence of our results also on this parameter is weak. Because of the truncated OPE, in general, the zeros of the sum rules for strong coupling constants and residues do not coincide. These points and close to these points are an artifact of the using truncated OPE and hence the "working" region of $\cos\theta$ should be far from these region. In Figs. 4–6 (as an example) we present the dependence of the coupling constants g_1, g_2 and g_3 of this transition on $\cos\theta$, where $\tan\theta = \beta$, at three fixed values of s_0 and at a fixed value of M^2. From these figures it is easily seen that the coupling constants g_1, g_2 and g_3 are practically unchanged while $\cos\theta$ is varying in the domain $-0.5 \le \cos\theta \le 0.3$ and weakly depend on s_0. Plotting all the considered strong coupling constants for all allowed transitions versus $\cos\theta$, we see that this working region is approximately common and optimal one to achieve reliable sum rules for all cases. From our analysis we obtain $g_1^{\Xi_c^{*+}\to\Xi_c^{'+}\rho^0} = (2.6 \pm 0.5) \times$ GeV^{-1}, $g_2^{\Xi_c^{*+}\to\Xi_c^{'+}\rho^0} = (0.9 \pm 0.2)$ GeV^{-2}, $g_3^{\Xi_c^{*+}\to\Xi_c^{'+}\rho^0} = (22 \pm 4)$ GeV^{-2}. The results for the coupling constants of other transitions are put into the Tables III and IV. For completeness, in these Tables we also present results of the nonrelativistic quark model (NRQM) on these couplings in terms of a constant c. From these tables we can conclude that predictions of the general and the Ioffe currents are very close to each other. We also see that ratios of the decays considered are also in good agreement with the predictions of the nonrelativistic quark model. Finally it should be noted that some of the coupling constants related with ρ^0, ω and ϕ mesons were studied in [23] with the Ioffe interpolating currents. The results obtained in that work do partially agree or disagree as compared to our predictions.

IV. CONCLUSION

In the present work, we have studied the $B_Q^*B_QV$ vertices within the LCSR method. These vertices are parametrized with three coupling constants. We have calculated them for all the $B_Q^*B_QV$ transitions with light vector mesons. The main result is that the correlation functions responsible for the coupling of the light vector mesons with the heavy sextet baryons of the spin-3/2 and the heavy sextet and antitriplet baryons of the spin-1/2 are described in terms of only one invariant function for each Lorenz structure while the relations between the different transitions are structure independent.

APPENDIX

In this Appendix we present the expressions of the correlation functions in terms of invariant function Π_1 involving ρ, K^*, ω and ϕ mesons.

(i) Correlation functions responsible for the sextet-sextet transitions:

$$\Pi^{\Sigma_b^{*0}\to\Sigma_b^0\rho^0} = \frac{1}{\sqrt{2}}\Big[\Pi_1(u, d, b) - \Pi_1(d, u, b)\Big],$$

$$\Pi^{\Sigma_b^{*+}\to\Sigma_b^+\rho^0} = \sqrt{2}\Pi_1(u, u, b),$$

$$\Pi^{\Sigma_b^{*-}\to\Sigma_b^-\rho^0} = -\sqrt{2}\Pi_1(d, d, b),$$

$$\Pi^{\Xi_b^{*0}\to\Xi_b^{'0}\rho^0} = \frac{1}{\sqrt{2}}\Pi_1(u, s, b),$$

$$\Pi^{\Xi_b^{*-}\to\Xi_b^{'-}\rho^0} = -\frac{1}{\sqrt{2}}\Pi_1(d, s, b),$$

$$\Pi^{\Sigma_b^{*+}\to\Sigma_b^0\rho^+} = \sqrt{2}\Pi_1(d, u, b),$$

$$\Pi^{\Sigma_b^{*0}\to\Sigma_b^-\rho^+} = \sqrt{2}\Pi_1(u, d, b),$$

$$\Pi^{\Xi_b^{*0}\to\Xi_b^{'-}\rho^+} = \Pi_1(d, s, b),$$

$$\Pi^{\Sigma_b^{*0}\to\Sigma_b^+\rho^-} = \sqrt{2}\Pi_1(d, u, b),$$

$$\Pi^{\Sigma_b^{*-}\to\Sigma_b^0\rho^-} = \sqrt{2}\Pi_1(u, d, b),$$

$$\Pi^{\Xi_b^{*-}\to\Xi_b^{'0}\rho^-} = \Pi_1(u, s, b),$$

$$\Pi^{\Xi_b^{*0}\to\Sigma_b^+K^{*-}} = \sqrt{2}\Pi_1(u, u, b),$$

$$\Pi^{\Xi_b^{*-}\to\Sigma_b^0K^{*-}} = \Pi_1(u, d, b),$$

$$\Pi^{\Omega_b^{*-}\to\Xi_b^{'0}K^{*-}} = \sqrt{2}\Pi_1(s, s, b),$$

$$\Pi^{\Sigma_b^{*+}\to\Xi_b^{'0}K^{*+}} = \sqrt{2}\Pi_1(u, u, b),$$

$$\Pi^{\Sigma_b^{*0}\to\Xi_b^{'-}K^{*+}} = \Pi_1(u, d, b),$$

$$\Pi^{\Xi_b^{*0}\to\Omega_b^-K^{*+}} = \sqrt{2}\Pi_1(s, s, b),$$

$$\Pi^{\Xi_b^{*0}\to\Sigma_b^0\bar{K}^{*0}} = \Pi_1(d, u, b),$$

$$\Pi^{\Xi_b^{*-}\to\Sigma_b^-\bar{K}^{*0}} = \sqrt{2}\Pi_1(d, d, b),$$

$$\Pi^{\Omega_b^{*-}\to\Xi_b^{'-}\bar{K}^{*0}} = \sqrt{2}\Pi_1(s, s, b),$$

$$\Pi^{\Sigma_b^{*0}\to\Xi_b^{'0}K^{*0}} = \Pi_1(d, u, b),$$

$$\Pi^{\Sigma_b^{*-}\to\Xi_b^{'-}K^{*0}} = \sqrt{2}\Pi_1(d, d, b),$$

$$\Pi^{\Xi_b^{*-}\to\Omega_b^-K^{*0}} = \sqrt{2}\Pi_1(s, s, b),$$

$$\Pi^{\Sigma_b^{*0}\to\Sigma_b^0\omega} = \frac{1}{\sqrt{2}}\Big[\Pi_1(u, d, b) + \Pi_1(d, u, b)\Big],$$

$$\Pi^{\Sigma_b^{*+}\to\Sigma_b^+\omega} = \sqrt{2}\Pi_1(u, u, b),$$

$$\Pi^{\Sigma_b^{*-}\to\Sigma_b^-\omega} = \sqrt{2}\Pi_1(d, d, b),$$

$$\Pi^{\Xi_b^{*0}\to\Xi_b^{'0}\omega} = \frac{1}{\sqrt{2}}\Pi_1(u, s, b),$$

$$\Pi^{\Xi_b^{*-}\to\Xi_b^{'-}\omega} = \frac{1}{\sqrt{2}}\Pi_1(d, s, b),$$

$$\Pi^{\Xi_b^{*0}\to\Xi_b^{'0}\phi} = \Pi_1(s, u, b),$$

$$\Pi^{\Xi_b^{*-}\to\Xi_b^{'-}\phi} = \Pi_1(s, d, b),$$

$$\Pi^{\Omega_b^{*-}\to\Omega_b^-\phi} = 2\Pi_1(s, s, b).$$

(ii) Correlation functions responsible for the sextet-antitriplet transitions:

$$\Pi^{\Xi_b^{*0}\to\Xi_b^0\rho^0} = \frac{1}{\sqrt{2}}\tilde{\Pi}_1(u,s,b),$$

$$\Pi^{\Xi_b^{*-}\to\Xi_b^-\rho^0} = \frac{1}{\sqrt{2}}\tilde{\Pi}_1(d,s,b),$$

$$\Pi^{\Sigma_b^{*0}\to\Lambda_b\rho^0} = \frac{1}{\sqrt{2}}\left[\tilde{\Pi}_1(u,d,b) + \tilde{\Pi}_1(d,u,b)\right],$$

$$\Pi^{\Sigma_b^{*-}\to\Lambda_b\rho^-} = \sqrt{2}\tilde{\Pi}_1(u,d,b),$$

$$\Pi^{\Xi_b^{*-}\to\Xi_b^-\rho^-} = \tilde{\Pi}_1(d,s,b),$$

$$\Pi^{\Sigma_b^{*+}\to\Lambda_b\rho^+} = -\sqrt{2}\tilde{\Pi}_1(d,u,b),$$

$$\Pi^{\Xi_b^{*0}\to\Xi_b^-\rho^+} = \tilde{\Pi}_1(u,s,b),$$

$$\Pi^{\Omega_b^{*-}\to\Xi_b^-\bar{K}^{*0}} = -\sqrt{2}\tilde{\Pi}_1(s,s,b),$$

$$\Pi^{\Xi_b^{*0}\to\Lambda_b\bar{K}^{*0}} = -\tilde{\Pi}_1(d,u,b),$$

$$\Pi^{\Sigma_b^{*0}\to\Xi_b^0 K^{*0}} = -\tilde{\Pi}_1(d,u,b),$$

$$\Pi^{\Sigma_b^{*-}\to\Xi_b^- K^{*0}} = \sqrt{2}\tilde{\Pi}_1(d,d,b),$$

$$\Pi^{\Sigma_b^{*0}\to\Xi_b^- K^{*+}} = \tilde{\Pi}_1(u,d,b),$$

$$\Pi^{\Omega_b^{*-}\to\Xi_b^0 K^{*-}} = \sqrt{2}\tilde{\Pi}_1(s,s,b),$$

$$\Pi^{\Xi_b^{*0}\to\Xi_b^0\omega} = \frac{1}{\sqrt{2}}\tilde{\Pi}_1(u,s,b),$$

$$\Pi^{\Xi_b^{*-}\to\Xi_b^-\omega} = \frac{1}{\sqrt{2}}\tilde{\Pi}_1(d,s,b),$$

$$\Pi^{\Sigma_b^{*0}\to\Lambda_b\omega} = \frac{1}{\sqrt{2}}\left[\tilde{\Pi}_1(u,d,b) - \tilde{\Pi}_1(d,u,b)\right],$$

$$\Pi^{\Xi_b^{*0}\to\Xi_b^0\phi} = -\tilde{\Pi}_1(s,u,b),$$

$$\Pi^{\Xi_b^{*-}\to\Xi_b^-\phi} = -\tilde{\Pi}_1(s,d,b).$$

The expressions for the charmed baryons can easily be obtained by making the replacement $b \to c$ and adding to charge of each baryon a positive unit charge.

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