

# Heavy baryon–light vector meson couplings in QCD

T. M. Aliev <sup>\*†</sup>, K. Azizi <sup>‡</sup>, M. Savcı <sup>§</sup>

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

<sup>‡</sup> Physics Division, Faculty of Arts and Sciences, Doğuş University,  
Acıbadem-Kadıköy, 34722 Istanbul, Turkey

## Abstract

The strong coupling constants of heavy baryons with light vector mesons are calculated in the framework of the light cone QCD sum rules using the most general form of the interpolating currents for the heavy baryons. It is shown that the sextet–sextet, sextet–antitriplet and antitriplet–antitriplet transitions are described by one invariant function for each class of transitions. The values of the electric and magnetic coupling constants for these transitions are obtained.

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

†permanent address: Institute of Physics, Baku, Azerbaijan

‡e-mail: kazizi@dogus.edu.tr

§e-mail: savci@metu.edu.tr

# 1 Introduction

Experimental and theoretical studies of charmed and bottom baryons are recognized to be one of the main research areas in particle physics. During the last few years quite interesting experimental observations have been obtained in heavy hadron spectroscopy. With the help of several refined measurements, the new states are discovered both in the charm and bottom sector (for recent experimental results see review [1]). These observations come from both the BaBar and BELLE, as well as from CDF and DØ Collaborations. LHC opens the possibilities for the discovery and detailed study of the new baryon states [2]. The considerable progress on the experimental side, has stimulated the theoretical investigation for understanding the dynamics of heavy flavor hadrons. Careful and comprehensive theoretical studies of the experimental results on heavy hadron spectroscopy and analysis of their weak and strong decays can provide essential knowledge on the quark structure of these hadrons. The strong coupling constants of the light pseudoscalar and vector mesons with heavy baryons are the main parameters for understanding the dynamics of heavy baryons. For this reason, reliable determination of the strong coupling constants of light pseudoscalar and vector mesons with heavy baryons within QCD receives special attention. Unfortunately, at the hadronic scale, QCD becomes nonperturbative and it makes impossible to calculate these coupling constants starting from QCD Lagrangian. Therefore, for calculation of these coupling constants some nonperturbative methods are needed. The QCD sum rule approach, which is based on the fundamental QCD Lagrangian, is one of the most attractive and applicable approaches [3]. In this work, we calculate the strong coupling constants of light vector mesons with heavy baryons within the light cone version of the QCD sum rules (LCSR) (for more about LCSR see [4]). The coupling constants of pseudoscalar mesons with heavy baryons is studied in detail in the same framework in [5].

The work is arranged in the following way. In section 2, the light cone sum rules for the coupling constants of light vector mesons with heavy baryons are obtained. In the following section, the numerical analysis of the obtained sum rules is performed and a comparison of our results with ones existing in literature is presented.

## 2 Light cone QCD sum rules for the heavy baryon–light vector meson couplings

In this section, we calculate the strong coupling constants of light–vector mesons with heavy baryons. Before starting to calculate these coupling constants few words about the  $SU(3)_f$  classification of the heavy baryons are in order. Baryons with a heavy single quark belong to either antisymmetric antitriplet  $\bar{3}_F$  or symmetric sextet  $6_F$  representations. Using the symmetry properties of the wave function, the spin of the light diquark is equal to zero for the antitriplet, while it is equal to one for the sextet. For this reason the total spin of the ground state baryons is 1/2 for  $\bar{3}_F$ , but it can be both 3/2 and 1/2 for  $6_F$ . In the present work we restrict ourselves by considering only spin 1/2 heavy baryons.

In order to obtain the strong coupling constants of light–vector mesons with heavy

baryons within the LCSR method, we consider the following correlation function:

$$\Pi^{(ij)} = i \int d^4x e^{ipx} \left\langle V(q) \left| \mathcal{T} \left\{ \eta_{B_2}^{(i)}(x) \bar{\eta}_{B_1}^{(j)}(0) \right\} \right| 0 \right\rangle, \quad (1)$$

where the indices  $i$  and  $j$  get two values and describe the sextet–sextet ( $i = 1, j = 1$ ), sextet–triplet ( $i = 1, j = 2$ ) and triplet–triplet ( $i = 2, j = 2$ ) transitions, respectively. In further respect, we introduce the definitions,  $\Pi^{(11)} = \Pi^{(1)}$ ,  $\Pi^{(12)} = \Pi^{(2)}$  and  $\Pi^{(22)} = \Pi^{(3)}$ , correspondingly. Note that, the  $V(q)$  in Eq. (1) denotes the light–vector mesons ( $\rho, \omega, K^*, \phi$ ) with momentum  $q$  and  $\eta$  is the interpolating current of the heavy baryon.

The correlation function (1) can be calculated in two different kinematical regions, namely, in terms of hadrons (phenomenological side), as well as in deep euclidean region when  $p^2 \rightarrow -\infty$ , in quark and gluon degrees of freedom using operator product expansion (OPE) (theoretical side). Equating both representations with the help of dispersion relations allows us to obtain the sum rules that will be used for calculation of the strong coupling constants of the light–vector mesons with heavy baryons.

The phenomenological part of the correlation function can be obtained by saturating it with hadrons which carry the same quantum numbers as their interpolating currents. Separating contribution of the ground state baryons we get,

$$\Pi^{(ij)} = \frac{\left\langle 0 \left| \eta_{B_2}^{(i)} \right| B_2(p) \right\rangle \left\langle B_2(p) V(q) \left| B_1(p+q) \right\rangle \left\langle B_1(p+q) \left| \bar{\eta}_{B_1}^{(j)} \right| 0 \right\rangle}{(p^2 - m_2^2) [(p+q)^2 - m_1^2]} + \dots, \quad (2)$$

where  $m_1$  and  $m_2$  are the masses of the initial and final baryons and dots represent contributions from higher states and continuum. The matrix elements appearing in Eq. (2) are defined as follows:

$$\left\langle 0 \left| \eta_B^{(i)} \right| B(p) \right\rangle = \lambda_i \bar{u}(p), \quad (3)$$

$$\left\langle B_2(p) V(q) \left| B_1(p+q) \right\rangle = \bar{u}(p) \left[ f_1 \gamma_\mu - i f_2 \sigma_{\mu\nu} q^\nu \frac{1}{m_1 + m_2} \right] u(p+q) \varepsilon^\mu, \quad (4)$$

where  $\lambda_i$  are the residues of the heavy baryons and  $\varepsilon^\mu$  are the momentum and vector polarization of the vector meson,  $f_1$  and  $f_2$  are the charge and magnetic form factors, respectively.

Performing summation over spins of the baryons and using Eqs. (2)–(4), we get for the phenomenological part,

$$\Pi^{(ij)} = i \frac{\lambda_i \lambda_j}{(p^2 - m_2^2) [(p+q)^2 - m_1^2]} \left\{ (f_1 + f_2) \not{p} \not{q} + 2 f_1 (\varepsilon \cdot p) \not{p} + \text{other structures} \right\}. \quad (5)$$

The reason why we choose the structures  $\not{p} \not{q}$  and  $(\varepsilon \cdot p) \not{p}$  is that they show the best convergence.

In calculating the theoretical part from the QCD side, explicit expressions of the interpolating currents of heavy baryons are needed. Using the fact that the sextet (antitriplet) current should be symmetric (antisymmetric) with respect to the light quarks, the most

general form of the interpolating currents for the spin-1/2 sextet and antitriplet baryons can be written in the following form,

$$\begin{aligned}\eta_Q^{(s)} &= -\frac{1}{\sqrt{2}}\epsilon^{abc}\left\{\left(q_1^{aT}CQ^b\right)\gamma_5q_2^c - \left(Q^{aT}Cq_2^b\right)\gamma_5q_1^c + \beta\left[\left(q_1^{aT}C\gamma_5Q^b\right)q_2^c - \left(Q^{aT}C\gamma_5q_2^b\right)q_1^c\right]\right\}, \\ \eta_Q^{(a)} &= \frac{1}{\sqrt{6}}\epsilon^{abc}\left\{2\left(q_1^{aT}Cq_2^b\right)\gamma_5Q^c + \left(q_1^{aT}CQ^b\right)\gamma_5q_2^c + \left(Q^{aT}Cq_2^b\right)\gamma_5q_1^c\right. \\ &\quad \left.+ \beta\left[2\left(q_1^{aT}C\gamma_5q_2^b\right)Q^c + \left(q_1^{aT}C\gamma_5Q^b\right)q_2^c + \left(Q^{aT}C\gamma_5q_2^b\right)q_1^c\right]\right\},\end{aligned}\tag{6}$$

where the superscripts  $s$  and  $a$  refer to the sextet and antitriplet, respectively, and subscripts  $a, b, c$  are the color indices,  $C$  is the charge conjugation operator,  $\beta$  is an arbitrary parameter, and  $\beta = -1$  case describes the Ioffe current. The light quark fields  $q_1$  and  $q_2$  for the sextet and antitriplet are given in Table 1.

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

Table 1: The light quarks  $q_1$  and  $q_2$  for the sextet and antitriplet baryons

Before calculating the correlation functions responsible for all sextet–sextet–vector mesons (SSV), sextet–antitriplet–vector mesons (SAV) and antitriplet–antitriplet–vector mesons (AAV) transitions from the QCD side, we find the relations among invariant functions for each of the above–mentioned classes (see also [5–10]). As a result, we find that the magnetic and electric couplings for each class of transitions are described in terms of only one invariant function. Of course, the invariant functions for each class of transitions are different from each other in the general case. It should also be noted that the relations among the invariant functions are all structure independent.

After these remarks, we can proceed to establish relations among invariant functions involving the couplings of sextet–sextet transitions. Let us first consider  $\Sigma_b^0 \rightarrow \Sigma_b^0\rho^0$  transition. The invariant function responsible for this transition can schematically be written as:

$$\Pi^{\Sigma_b^0 \rightarrow \Sigma_b^0\rho^0} = g_{\rho\bar{u}u}\Pi_1^{(1)}(u, d, b) + g_{\rho\bar{d}d}\Pi_1^{\prime(1)}(u, d, b) + g_{\rho\bar{b}b}\Pi_2^{(1)}(u, d, b),\tag{7}$$

where we have introduced the formal notations,

$$\begin{aligned}
\Pi_1^{(1)}(u, d, b) &= \langle \bar{u}u | \Sigma_b^0 \bar{\Sigma}_b^0 | 0 \rangle , \\
\Pi_2^{(1)}(u, d, b) &= \langle \bar{b}b | \Sigma_b^0 \bar{\Sigma}_b^0 | 0 \rangle , \\
\Pi_1^{\prime(1)}(u, d, b) &= \langle \bar{d}d | \Sigma_b^0 \bar{\Sigma}_b^0 | 0 \rangle .
\end{aligned} \tag{8}$$

The interpolating current of  $\rho^0$  is formally written in the form,

$$J_\mu^{\rho^0} = \sum_{u,d,b} g_{\rho\bar{q}q} \bar{q} \gamma_\mu q , \tag{9}$$

where we have set  $g_{\rho^0\bar{b}b} = 0$  and  $g_{\rho^0\bar{u}u} = -g_{\rho^0\bar{d}d} = 1/\sqrt{2}$ . Physically, each term on the right hand side of Eq. (7) describes emission of the  $\rho^0$  meson from  $u$ ,  $d$  and  $b$  quarks of the  $\Sigma_b^0$  baryon, respectively. Since interpolating current of  $\Sigma_b^0$  is symmetric under the exchange of  $u$  and  $d$  quarks, it leads to the result,  $\Pi_1^{(1)}(u, d, b) = \Pi_1^{(1)}(d, u, b)$ . Using the definitions given in Eq. (8) and taking the remark after Eq. (9) into consideration, we obtain,

$$\Pi^{\Sigma_b^0 \rightarrow \Sigma_b^0 \rho^0} = \frac{1}{\sqrt{2}} \left[ \Pi_1^{(1)}(u, d, b) - \Pi_1^{(1)}(d, u, b) \right] . \tag{10}$$

Obviously, in the  $SU(2)_f$  limit  $\Pi^{\Sigma_b^0 \rightarrow \Sigma_b^0 \rho^0} = 0$ .

The invariant function describing  $\Sigma_b^+ \rightarrow \Sigma_b^+ \rho^0$  ( $\Sigma_b^- \rightarrow \Sigma_b^- \rho^0$ ) can be obtained from the invariant function for the  $\Sigma_b^0 \rightarrow \Sigma_b^0 \rho^0$  transition with the help of the replacement,  $d \rightarrow u$  ( $u \rightarrow d$ ), and using the fact that  $\Sigma_b^0 = -\sqrt{2}\Sigma_b^+$  ( $\sqrt{2}\Sigma_b^-$ ). As a result, we obtain,

$$4\Pi_1^{(1)}(u, u, b) = -2 \langle \bar{u}u | \Sigma_b^+ \bar{\Sigma}_b^+ | 0 \rangle , \tag{11}$$

$$4\Pi_1^{(1)}(d, d, b) = 2 \langle \bar{d}d | \Sigma_b^- \bar{\Sigma}_b^- | 0 \rangle . \tag{12}$$

Since  $\Sigma_b^{+(-)}$  contains two  $u(d)$  quarks, there are four possible ways for emitting  $\rho^0$  from the  $u(d)$  quark. Hence, the related invarian functions are obtained as:

$$\Pi^{\Sigma_b^+ \rightarrow \Sigma_b^+ \rho^0} = \sqrt{2}\Pi_1^{(1)}(u, u, b) , \tag{13}$$

$$\Pi^{\Sigma_b^- \rightarrow \Sigma_b^- \rho^0} = \sqrt{2}\Pi_1^{(1)}(d, d, b) . \tag{14}$$

The result for the invariant function responsible for the  $\Xi_b^{\prime-(0)} \rightarrow \Xi_b^{\prime-(0)} \rho^0$  transitions can easily be obtained from the result for the  $\Sigma_b^0 \rightarrow \Sigma_b^0 \rho^0$  transition using the fact that  $\Xi_b^{\prime 0} = \Sigma_b^0(d \rightarrow s)$  and  $\Xi_b^{\prime -} = \Sigma_b^0(u \rightarrow s)$ , i.e.,

$$\Pi^{\Xi_b^{\prime 0} \rightarrow \Xi_b^{\prime 0} \rho^0} = \frac{1}{\sqrt{2}} \Pi_1^{(1)}(u, s, b) , \tag{15}$$

$$\Pi^{\Xi_b^{\prime -} \rightarrow \Xi_b^{\prime -} \rho^0} = -\frac{1}{\sqrt{2}} \Pi_1^{(1)}(d, s, b) . \tag{16}$$

How can one find the relations among the invariant functions in the presence of the charged  $\rho^\pm$  mesons? In order to answer this question, we again consider the matrix element,  $\langle \bar{d}d | \Sigma_b^0 \bar{\Sigma}_b^0 | 0 \rangle$ . This matrix element means that the  $d$  quarks from  $\Sigma_b^0$  and  $\bar{\Sigma}_b^0$  form the final

$\bar{d}d$  state, while  $u$  and  $b$  quarks are the spectators. The matrix element  $\langle \bar{u}d | \Sigma_b^+ \bar{\Sigma}_b^0 | 0 \rangle$  corresponds to the case when  $d$  quark from  $\bar{\Sigma}_b^0$  and  $u$  quark from  $\Sigma_b^+$  form the  $\bar{u}d$  state, again the remaining  $u$  and  $b$  quarks being the spectators. This fact allows us to comment that these matrix elements are proportional to each other. A detailed calculation shows that, these matrix elements are related to each other through,

$$\begin{aligned} \Pi^{\Sigma_b^0 \rightarrow \Sigma_b^+ \rho^-} &= \langle \bar{u}d | \Sigma_b^+ \bar{\Sigma}_b^0 | 0 \rangle = -\sqrt{2} \langle \bar{d}d | \Sigma_b^0 \bar{\Sigma}_b^0 | 0 \rangle \\ &= -\sqrt{2} \Pi_1^{(1)}(d, u, b) . \end{aligned} \quad (17)$$

Making the replacement  $u \leftrightarrow d$ , from Eq. (17) we get,

$$\Pi^{\Sigma_b^0 \rightarrow \Sigma_b^- \rho^+} = \sqrt{2} \Pi_1^{(1)}(u, d, b) . \quad (18)$$

In calculating the coupling constants of SSV, SAV and AAV, it is enough to consider the transitions  $\Sigma_b^0 \rightarrow \Sigma_b^0 V$ ,  $\Xi_b^{\prime 0} \rightarrow \Xi_b^{\prime 0} V$  and  $\Xi_b^0 \rightarrow \Xi_b^0 V$ , respectively, since all other strong couplings can be achieved from these results with the help of corresponding replacements among the quarks. In order to obtain the correlation functions in terms of  $\Pi_1$ , which describe transitions among sextet–sextet, sextet–antitriplet and antitriplet–antitriplet with other vector mesons, similar calculations can be done. These results are presented in Appendix A.

We can now proceed to calculate the invariant functions  $\Pi_1^{(i)}$  ( $i = 1, 2, 3$ ) responsible for  $\Sigma_b^0 \rightarrow \Sigma_b^0 \rho^0$ ,  $\Xi_b^{\prime 0} \rightarrow \Xi_b^{\prime 0} \rho^0$  and  $\Xi_b^0 \rightarrow \Xi_b^0 \rho^0$  transitions from QCD side. These invariant functions can be calculated in deep Euclidian region  $-p^2 \rightarrow \infty$ ,  $-(p+q)^2 \rightarrow \infty$  using the operator product expansion. The main nonperturbative ingredient in the calculations are the distribution amplitude (DA's) of the vector mesons. These distribution amplitudes appear in determination of the matrix elements of nonlocal operators  $\langle V(q) | \bar{q}(x) \Gamma q(0) | 0 \rangle$  and  $\langle V(q) | \bar{q}(x) G_{\mu\nu} q(0) | 0 \rangle$ , where  $\Gamma$  is any Dirac matrix. Up to twist twist-4 accuracy, the expressions for the distribution functions of vector mesons can be found in [11, 12].

In calculating the correlation functions responsible for afore-mentioned decays from QCD side, the expressions of light and heavy quark propagators are also needed. The light quark propagator in an external field is calculated in [13] whose expression is given as

$$\begin{aligned} S_q(x) &= \frac{i\not{x}}{2\pi^2 x^4} - \frac{m_q}{4\pi^2 x^2} - \frac{\langle \bar{q}q \rangle}{12} \left( 1 - i \frac{m_q}{4} \not{x} \right) - \frac{x^2}{192} m_0^2 \langle \bar{q}q \rangle \left( 1 - i \frac{m_q}{6} \not{x} \right) \\ &\quad - i g_s \int_0^1 du \left[ \frac{\not{x}}{16\pi^2 x^2} G_{\mu\nu}(ux) \sigma_{\mu\nu} - \frac{i}{4\pi^2 x^2} u x^\mu G_{\mu\nu}(ux) \gamma^\nu \right. \\ &\quad \left. - i \frac{m_q}{32\pi^2} G_{\mu\nu} \sigma^{\mu\nu} \left( \ln \left( \frac{-x^2 \Lambda^2}{4} \right) + 2\gamma_E \right) \right] , \end{aligned} \quad (19)$$

where  $\gamma_E \simeq 0.577$  is the Euler Constant, and  $\Lambda$  is the scale parameter and it is chosen as the factorization scale  $\Lambda = (0.5 \div 1) \text{ GeV}$  (for more detail see [14]).

The propagator for the heavy quark is [15]:

$$\begin{aligned} S_Q(x) &= \frac{m_Q^2}{4\pi^2} \frac{K_1(m_Q \sqrt{-x^2})}{\sqrt{-x^2}} - i \frac{m_Q^2 \not{x}}{4\pi^2 x^2} K_2(m_Q \sqrt{-x^2}) \\ &\quad - i g_s \int \frac{d^4 k}{(2\pi)^4} e^{-ikx} \int_0^1 du \left[ \frac{\not{k} + m_Q}{2(m_Q^2 - k^2)^2} G^{\mu\nu}(ux) \sigma_{\mu\nu} + \frac{u}{m_Q^2 - k^2} x_\mu G^{\mu\nu} \gamma_\nu \right] , \end{aligned} \quad (20)$$

where  $K_i$  are the modified Bessel function of the second kind.

Using the expressions of light and heavy quark propagators, as well as definitions of the DA's for the vector mesons, and after lengthy calculations one can obtain the correlation function from QCD part.

As has already been noted, the relations among the correlation functions for the considered transitions are structure independent, but their explicit expressions are structure dependent. For this reason we introduce new indices  $\alpha$  in the correlation function, where  $\alpha = 1$  stands for the choice of the structure  $(\varepsilon \cdot p)\not{p}$  and  $\alpha = 2$  for the structure  $\not{p}\not{q}$ .

Equating the coefficients of the structures,  $(\varepsilon p)\not{p}$  and  $\not{p}\not{q}$  for the hadronic and QCD sides and performing Borel transformations over the variables  $p^2$  and  $(p+q)^2$ , which suppress the contributions of the continuum and higher states, we finally get the sum rules for the strong coupling constants of light vector mesons with sextet and antitriplet heavy baryons as:

$$f_\alpha^{(i)} = \frac{1}{\lambda_1^{(i)} \lambda_2^{(i)}} e^{\frac{m_1^{(i)2}}{M_1^2} + \frac{m_2^{(i)2}}{M_2^2} + \frac{m_V^2}{M_1^2 + M_2^2}} \Pi_\alpha^{(i)}, \quad (21)$$

where  $M_1^2$  and  $M_2^2$  are the Borel parameters corresponding to the initial and final heavy baryons, respectively. In the problem under consideration the masses of the initial and final heavy baryons are very close to each other, hence we can take  $M_1^2 = M_2^2 \equiv 2M^2$ . Note that the residues of the sextet and antitriplet heavy baryons are calculated in [16]. This leads to the result that for numerical calculation of the sum rules, the DA's are needed to be evaluated only at  $u_0 = \frac{M_1^2}{M_1^2 + M_2^2} = \frac{1}{2}$ .

### 3 Numerical analysis

Having already obtained the sum rules for the strong coupling constants of light vector mesons with heavy baryons, we can now proceed evaluating them numerically. The essential ingredients of the LCSR are the DA's of the light vector mesons. Explicit expressions of DA's for the vector mesons and the values of the parameters entering to the expressions of DA's are given in [11, 12]. The residues of the heavy baryons are calculated in [16].

In addition to the DA's and other input parameters, the sum rules for SSV, SAV and AAV transitions contain three auxiliary parameters : the continuum threshold  $s_0$ , Borel parameter  $M^2$  and the parameter  $\beta$  entering the expressions of the interpolating current. For this reason we try to find such regions of these parameters where strong coupling constants are practically independent of them.

In finding the working region of  $M^2$ , we require that the continuum and higher state contributions should be less than half of the dispersion integral, and additionally the contribution of the higher terms with the power  $1/M^2$  be 25% less than the total result. These two restrictions lead to the result that the "working region" of  $M^2$  is  $15 \text{ GeV}^2 \leq M^2 \leq 30 \text{ GeV}^2$  (for the bottom baryons) and  $4 \text{ GeV}^2 \leq M^2 \leq 12 \text{ GeV}^2$  (for the charmed ones). The continuum threshold is varied in the region  $(m_B + 0.5)^2 \text{ GeV}^2 \leq s_0 \leq (m_B + 0.7)^2 \text{ GeV}^2$ . To be more illustrative about how the numerical analysis is performed, as an example, we consider the  $\Xi_b^{\prime 0} \rightarrow \Xi_b^{\prime 0} \rho^0$  transition. In Figs. (1) and (2) we present the dependence of  $f_1$  and  $f_2$  on  $M^2$  for the above-mentioned transition, at five different values of  $\beta$ , and at a

$f_1$ channel	Bottom Baryons		$f_1$ channel	Charmed Baryons	
	General current	Ioffe current		General current	Ioffe current
$f_1^{\Xi_b^{\prime 0} \rightarrow \Xi_b^{\prime 0} \rho^0}$	$2.2 \pm 0.7$	$2.0 \pm 0.7$	$f_1^{\Xi_c^{\prime +} \rightarrow \Xi_c^{\prime +} \rho^0}$	$2.5 \pm 0.8$	$5.0 \pm 1.7$
$f_1^{\Sigma_b^0 \rightarrow \Sigma_b^- \rho^+}$	$4.5 \pm 1.5$	$3.4 \pm 1.1$	$f_1^{\Sigma_c^+ \rightarrow \Sigma_c^0 \rho^+}$	$4.0 \pm 1.3$	$3.4 \pm 1.1$
$f_1^{\Xi_b^{\prime 0} \rightarrow \Sigma_b^+ K^{*-}}$	$6.0 \pm 2.0$	$3.9 \pm 1.3$	$f_1^{\Xi_c^{\prime +} \rightarrow \Sigma_c^{++} K^{*-}}$	$5.0 \pm 1.7$	$3.8 \pm 1.3$
$f_1^{\Omega_b^- \rightarrow \Xi_b^{\prime 0} \bar{K}^{*-}}$	$6.0 \pm 2.0$	$4.8 \pm 1.6$	$f_1^{\Omega_c^0 \rightarrow \Xi_c^{\prime +} \bar{K}^{*-}}$	$7.0 \pm 2.0$	$14.0 \pm 5.0$
$f_1^{\Sigma_b^+ \rightarrow \Sigma_b^+ \omega}$	$4.0 \pm 1.3$	$3.0 \pm 1.0$	$f_1^{\Sigma_c^{++} \rightarrow \Sigma_c^{++} \omega}$	$3.5 \pm 1.2$	$3.0 \pm 1.0$
$f_1^{\Xi_b^{\prime 0} \rightarrow \Xi_b^{\prime 0} \omega}$	$2.1 \pm 0.7$	$1.7 \pm 0.6$	$f_1^{\Xi_c^{\prime +} \rightarrow \Xi_c^{\prime +} \omega}$	$2.4 \pm 0.8$	$4.9 \pm 1.6$
$f_1^{\Xi_b^{\prime 0} \rightarrow \Xi_b^{\prime 0} \phi}$	$5.0 \pm 1.7$	$2.6 \pm 0.9$	$f_1^{\Xi_c^{\prime +} \rightarrow \Xi_c^{\prime +} \phi}$	$4.0 \pm 1.3$	$2.5 \pm 0.8$
$f_1^{\Omega_b^- \rightarrow \Omega_b^- \phi}$	$10.0 \pm 3.4$	$7.0 \pm 2.4$	$f_1^{\Omega_c^0 \rightarrow \Omega_c^0 \phi}$	$11.0 \pm 4.0$	$23.0 \pm 8.0$

Table 2: The values of the strong coupling constants  $f_1$  for the transitions among the sextet–sextet heavy baryons with vector mesons.

fixed value of  $s_0$ . From these figures, one can conclude that the strong coupling constants,  $f_1$  and  $f_2$  are practically independent of  $M^2$  when it varies in its own ‘working region’.

In Figs. (3) and (4) we present the dependence of  $f_1$  and  $f_2$  on  $\cos\theta$ , at three fixed values of  $s_0$  and  $M^2$ , where  $\tan\theta = \beta$ . We observe from these figures that when  $\cos\theta$  is varied in the region  $-0.5 \leq \cos\theta \leq 0.3$ , the results are insensitive to the variation of  $\beta$ . From these figures, we obtain,  $f_1^{\Xi_b^{\prime 0} \rightarrow \Xi_b^{\prime 0} \rho^0} = 2.2 \pm 0.7$  and  $f_2^{\Xi_b^{\prime 0} \rightarrow \Xi_b^{\prime 0} \rho^0} = 30 \pm 10$ . Similar analysis for the other couplings of vector mesons are carried out and the results for  $f_1$  and  $f_2$  are presented in Tables (2)–(7), respectively. The errors presented in these Tables are due to the variation of the auxiliary parameters, as well as uncertainties in the values of the input parameters.

The results on the coupling constants of the light vector mesons with heavy baryons presented in Tables (2)–(7) lead to the following conclusions:

- For the coupling constant,  $f_1$ : The predictions for the strong coupling constant  $f_1$  which are obtained using the most general and Ioffe currents ( $\beta = -1$ ) disagree considerably from each other, especially for the channels,  $\Xi_b^{\prime 0} \rightarrow \Xi_b^{\prime 0} \phi$ ,  $\Xi_c^{\prime +} \rightarrow \Xi_c^{\prime +} \omega$ ,  $\Omega_c^0 \rightarrow \Xi_c^{\prime +} K^{*-}$ ,  $\Omega_c^0 \rightarrow \Omega_c^0 \phi$ ,  $\Xi_c^+ \rightarrow \Xi_c^+ \rho^0$ ,  $\Xi_c^{\prime +} \rightarrow \Xi_c^+ \rho^0$ ,  $\Xi_c^{\prime +} \rightarrow \Xi_c^{\prime +} \rho^0$ ,  $\Xi_b^{\prime 0} \rightarrow \Xi_b^0 \rho^0$ ,  $\Xi_c^+ \rightarrow \Xi_c^+ \omega$ ,  $\Xi_b^{\prime 0} \rightarrow \Xi_b^- \omega$ ,  $\Sigma_b^- \rightarrow \Lambda_b^0 \rho^-$  and  $\Sigma_c^0 \rightarrow \Lambda_c^+ \rho^-$ .
- For the  $f_2$  channel: predictions of the general and Ioffe currents for the strong coupling constants of SSV and AAV show considerable differences. These discrepancies can be attributed to the fact that,  $\beta = -1$  lies outside the stability region of  $\beta$  and it makes the predictions less reliable.

Our final remark in this section is that the coupling constants of the  $\Xi_c^0 \rightarrow \Xi_c^0 \rho^0$ ,  $\Xi_c^0 \rightarrow \Xi_c^0 \omega$  and  $\Xi_c^0 \rightarrow \Xi_c^0 \phi$  transitions are also calculated in the HQET [17] within



$f_1$ channel	Bottom Baryons		$f_1$ channel	Charmed Baryons	
	General current	Ioffe current		General current	Ioffe current
$f_1^{\Xi_b^{\prime 0} \rightarrow \Xi_b^0 \rho^0}$	$1.4 \pm 0.5$	$0.6 \pm 0.2$	$f_1^{\Xi_c^{\prime +} \rightarrow \Xi_c^+ \rho^0}$	$1.5 \pm 0.5$	$0.7 \pm 0.2$
$f_1^{\Xi_b^{\prime 0} \rightarrow \Xi_b^- K^{*+}}$	$2.5 \pm 0.8$	$1.3 \pm 0.4$	$f_1^{\Xi_c^{\prime +} \rightarrow \Xi_c^0 K^{*+}}$	$2.3 \pm 0.8$	$1.2 \pm 0.4$
$f_1^{\Sigma_b^- \rightarrow \Lambda_b^0 \rho^-}$	$2.8 \pm 0.9$	$0.8 \pm 0.3$	$f_1^{\Sigma_c^0 \rightarrow \Lambda_c^+ \rho^-}$	$2.6 \pm 0.9$	$0.6 \pm 0.2$
$f_1^{\Sigma_b^0 \rightarrow \Xi_b^0 \bar{K}^{*0}}$	$2.6 \pm 0.8$	$1.5 \pm 0.5$	$f_1^{\Sigma_c^+ \rightarrow \Xi_c^+ \bar{K}^{*0}}$	$2.2 \pm 0.7$	$0.8 \pm 0.3$
$f_1^{\Omega_b^- \rightarrow \Xi_b^- \bar{K}^{*0}}$	$3.5 \pm 1.2$	$2.0 \pm 0.6$	$f_1^{\Omega_c^0 \rightarrow \Xi_c^0 \bar{K}^{*0}}$	$3.3 \pm 1.1$	$1.7 \pm 0.6$
$f_1^{\Xi_b^{\prime 0} \rightarrow \Xi_b^- \omega}$	$1.3 \pm 0.4$	$0.5 \pm 0.2$	$f_1^{\Xi_c^{\prime +} \rightarrow \Xi_c^0 \omega}$	$1.2 \pm 0.4$	$1.1 \pm 0.4$
$f_1^{\Xi_b^{\prime 0} \rightarrow \Xi_b^0 \phi}$	$2.6 \pm 0.9$	$2.0 \pm 0.7$	$f_1^{\Xi_c^{\prime +} \rightarrow \Xi_c^+ \phi}$	$2.1 \pm 0.7$	$1.4 \pm 0.5$

Table 3: The values of the strong coupling constants  $f_1$  for the transitions among the sextet–antitriplet heavy baryons with vector mesons.

$f_1$ channel	Bottom Baryons		$f_1$ channel	Charmed Baryons	
	General current	Ioffe current		General current	Ioffe current
$f_1^{\Xi_b^0 \rightarrow \Xi_b^0 \rho^0}$	$3.1 \pm 1.1$	$2.5 \pm 0.8$	$f_1^{\Xi_c^+ \rightarrow \Xi_c^+ \rho^0}$	$6.0 \pm 2.0$	$1.5 \pm 0.5$
$f_1^{\Xi_b^- \rightarrow \Lambda_b^0 K^{*-}}$	$5.0 \pm 1.7$	$4.6 \pm 1.5$	$f_1^{\Xi_c^0 \rightarrow \Lambda_c^+ K^{*-}}$	$4.6 \pm 1.5$	$4.1 \pm 1.4$
$f_1^{\Xi_b^0 \rightarrow \Xi_b^0 \omega}$	$2.8 \pm 0.9$	$2.3 \pm 0.8$	$f_1^{\Xi_c^+ \rightarrow \Xi_c^+ \omega}$	$5.5 \pm 1.8$	$1.2 \pm 0.4$
$f_1^{\Lambda_b^0 \rightarrow \Lambda_b^0 \omega}$	$5.2 \pm 1.7$	$4.6 \pm 1.5$	$f_1^{\Lambda_c^+ \rightarrow \Lambda_c^+ \omega}$	$4.9 \pm 1.6$	$4.3 \pm 1.4$
$f_1^{\Lambda_b^0 \rightarrow \Lambda_b^0 \phi}$	$5.0 \pm 1.7$	$4.7 \pm 1.6$	$f_1^{\Lambda_c^+ \rightarrow \Lambda_c^+ \phi}$	$4.6 \pm 1.5$	$4.1 \pm 1.4$

Table 4: The values of the strong coupling constants  $f_1$  for the transitions among the antitriplet–antitriplet heavy baryons with vector mesons.

the same framework as our work. Our predictions on the coupling constants of these transitions and those given in [17] are in a good agreement with each other.

## 4 Acknowledgment

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$f_2$ channel	Bottom Baryons		$f_2$ channel	Charmed Baryons	
	General current	Ioffe current		General current	Ioffe current
$f_2^{\Xi_b^{\prime 0} \rightarrow \Xi_b^{\prime 0} \rho^0}$	$30.0 \pm 10.0$	$39.0 \pm 13.0$	$f_2^{\Xi_c^{\prime +} \rightarrow \Xi_c^{\prime +} \rho^0}$	$16.0 \pm 5.2$	$18.0 \pm 6.0$
$f_2^{\Sigma_b^0 \rightarrow \Sigma_b^- \rho^+}$	$55.0 \pm 18.0$	$72.0 \pm 24.0$	$f_2^{\Sigma_c^+ \rightarrow \Sigma_c^0 \rho^+}$	$27.0 \pm 9.0$	$33.0 \pm 11.0$
$f_2^{\Xi_b^{\prime 0} \rightarrow \Sigma_b^+ K^{*-}}$	$60.0 \pm 20.0$	$80.0 \pm 27.0$	$f_2^{\Xi_c^{\prime +} \rightarrow \Sigma_c^{++} K^{*-}}$	$30.0 \pm 10.0$	$36.0 \pm 12.0$
$f_2^{\Omega_b^- \rightarrow \Xi_b^{\prime 0} \bar{K}^{*-}}$	$70.0 \pm 23.0$	$88.0 \pm 29.0$	$f_2^{\Omega_c^0 \rightarrow \Xi_c^{\prime +} \bar{K}^{*-}}$	$35.0 \pm 12.0$	$41.0 \pm 14.0$
$f_2^{\Sigma_b^+ \rightarrow \Sigma_b^+ \omega}$	$50.0 \pm 17.0$	$64.0 \pm 21.0$	$f_2^{\Sigma_c^{++} \rightarrow \Sigma_c^{++} \omega}$	$24.0 \pm 8.0$	$29.0 \pm 9.5$
$f_2^{\Xi_b^{\prime 0} \rightarrow \Xi_b^{\prime 0} \omega}$	$27.0 \pm 9.0$	$34.0 \pm 11.0$	$f_2^{\Xi_c^{\prime +} \rightarrow \Xi_c^{\prime +} \omega}$	$15.0 \pm 5.0$	$16.0 \pm 6.3$
$f_2^{\Xi_b^{\prime 0} \rightarrow \Xi_b^{\prime 0} \phi}$	$45.0 \pm 15.0$	$57.0 \pm 19.0$	$f_2^{\Xi_c^{\prime +} \rightarrow \Xi_c^{\prime +} \phi}$	$21.0 \pm 7.0$	$26.0 \pm 8.6$
$f_2^{\Omega_b^- \rightarrow \Omega_b^- \phi}$	$95.0 \pm 32.0$	$125.0 \pm 42.0$	$f_2^{\Omega_c^0 \rightarrow \Omega_c^0 \phi}$	$52.0 \pm 17.0$	$60.0 \pm 20.0$

Table 5: The values of the strong coupling constants  $f_2$  for the transitions among the sextet–sextet heavy baryons with vector mesons.

## Appendix A :

Here in this appendix we present the expressions of the correlation functions in terms of invariant function  $\Pi_1^{(i)}$  involving  $\rho$ ,  $K$ , *omega* and  $\phi$  mesons.

- Correlation functions responsible for the sextet–sextet transitions.

$$\begin{aligned}
\Pi^{\Sigma_b^+ \rightarrow \Sigma_b^0 \rho^+} &= \sqrt{2} \Pi_1^{(1)}(d, u, b) , \\
\Pi^{\Sigma_b^0 \rightarrow \Sigma_b^- \rho^+} &= \sqrt{2} \Pi_1^{(1)}(u, d, b) , \\
\Pi^{\Xi_b^{\prime 0} \rightarrow \Xi_b^{\prime -} \rho^+} &= \Pi_1^{(1)}(d, s, b) , \\
\Pi^{\Sigma_b^0 \rightarrow \Sigma_b^+ \rho^-} &= \sqrt{2} \Pi_1^{(1)}(d, u, b) , \\
\Pi^{\Sigma_b^- \rightarrow \Sigma_b^0 \rho^-} &= \sqrt{2} \Pi_1^{(1)}(u, d, b) , \\
\Pi^{\Xi_b^{\prime -} \rightarrow \Xi_b^{\prime 0} \rho^-} &= \Pi_1^{(1)}(u, s, b) , \\
\Pi^{\Xi_b^{\prime 0} \rightarrow \Sigma_b^+ K^{*-}} &= \sqrt{2} \Pi_1^{(1)}(u, u, b) , \\
\Pi^{\Xi_b^{\prime -} \rightarrow \Sigma_b^0 K^{*-}} &= \Pi_1^{(1)}(u, d, b) , \\
\Pi^{\Omega_b^- \rightarrow \Xi_b^{\prime 0} K^{*-}} &= \sqrt{2} \Pi_1^{(1)}(s, s, b) , \\
\Pi^{\Sigma_b^+ \rightarrow \Xi_b^{\prime 0} K^{*+}} &= \sqrt{2} \Pi_1^{(1)}(u, u, b) , \\
\Pi^{\Sigma_b^0 \rightarrow \Xi_b^{\prime -} K^{*+}} &= \Pi_1^{(1)}(u, d, b) , \\
\Pi^{\Xi_b^{\prime 0} \rightarrow \Omega_b^- K^{*+}} &= \sqrt{2} \Pi_1^{(1)}(s, s, b) ,
\end{aligned}$$

$f_2$ channel	Bottom Baryons		$f_2$ channel	Charmed Baryons	
	General current	Ioffe current		General current	Ioffe current
$f_2^{\Xi_b^{\prime 0} \rightarrow \Xi_b^0 \rho^0}$	$22.0 \pm 7.0$	$23.0 \pm 7.0$	$f_2^{\Xi_c^{\prime +} \rightarrow \Xi_c^+ \rho^0}$	$11.0 \pm 3.8$	$11.0 \pm 3.8$
$f_2^{\Xi_b^{\prime 0} \rightarrow \Xi_b^- K^{*+}}$	$34.0 \pm 11.0$	$38.0 \pm 13.0$	$f_2^{\Xi_c^{\prime +} \rightarrow \Xi_c^0 K^{*+}}$	$15.0 \pm 5.0$	$18.0 \pm 6.0$
$f_2^{\Sigma_b^- \rightarrow \Lambda_b^0 \rho^-}$	$40.0 \pm 13.0$	$42.0 \pm 14.0$	$f_2^{\Sigma_c^0 \rightarrow \Lambda_c^+ \rho^-}$	$16.0 \pm 5.3$	$18.0 \pm 6.0$
$f_2^{\Sigma_b^0 \rightarrow \Xi_b^0 \bar{K}^{*0}}$	$32.0 \pm 11.0$	$35.0 \pm 12.0$	$f_2^{\Sigma_c^+ \rightarrow \Xi_c^+ \bar{K}^{*0}}$	$13.0 \pm 4.3$	$15.0 \pm 5.0$
$f_2^{\Omega_b^- \rightarrow \Xi_b^- \bar{K}^0}$	$50.0 \pm 17.0$	$53.0 \pm 18.0$	$f_2^{\Omega_c^0 \rightarrow \Xi_c^0 \bar{K}^0}$	$20.0 \pm 7.0$	$26.0 \pm 8.5$
$f_2^{\Xi_b^{\prime 0} \rightarrow \Xi_b^- \omega}$	$20.3 \pm 7.0$	$20.0 \pm 7.0$	$f_2^{\Xi_c^{\prime +} \rightarrow \Xi_c^0 \omega}$	$8.0 \pm 2.7$	$10.0 \pm 3.0$
$f_2^{\Xi_b^{\prime 0} \rightarrow \Xi_b^0 \phi}$	$30.0 \pm 10.0$	$34.0 \pm 11.0$	$f_2^{\Xi_c^{\prime +} \rightarrow \Xi_c^+ \phi}$	$13.0 \pm 4.3$	$15.0 \pm 5.0$

Table 6: The values of the strong coupling constants  $f_2$  for the transitions among the sextet–antitriplet heavy baryons with vector mesons.

$f_2$ channel	Bottom Baryons		$f_2$ channel	Charmed Baryons	
	General current	Ioffe current		General current	Ioffe current
$f_2^{\Xi_b^0 \rightarrow \Xi_b^0 \rho^0}$	$5.0 \pm 1.7$	$12.0 \pm 3.8$	$f_2^{\Xi_c^+ \rightarrow \Xi_c^+ \rho^0}$	$7.5 \pm 2.5$	$5.7 \pm 1.9$
$f_2^{\Xi_b^- \rightarrow \Lambda_b^0 K^{*-}}$	$7.0 \pm 2.0$	$24.0 \pm 8.0$	$f_2^{\Xi_c^0 \rightarrow \Lambda_c^+ K^{*-}}$	$6.0 \pm 2.0$	$12.0 \pm 4.0$
$f_2^{\Xi_b^0 \rightarrow \Xi_b^0 \omega}$	$4.0 \pm 1.3$	$11.0 \pm 3.6$	$f_2^{\Xi_c^+ \rightarrow \Xi_c^+ \omega}$	$7.5 \pm 2.5$	$5.1 \pm 1.7$
$f_2^{\Lambda_b^0 \rightarrow \Lambda_b^0 \omega}$	$8.0 \pm 2.7$	$25.0 \pm 8.0$	$f_2^{\Lambda_c^+ \rightarrow \Lambda_c^+ \omega}$	$6.0 \pm 2.0$	$14.0 \pm 5.0$
$f_2^{\Lambda_b^0 \rightarrow \Lambda_b^0 \phi}$	$8.0 \pm 2.7$	$23.0 \pm 7.5$	$f_2^{\Lambda_c^+ \rightarrow \Lambda_c^+ \phi}$	$6.0 \pm 2.0$	$12.0 \pm 4.0$

Table 7: The values of the strong coupling constants  $f_2$  for the transitions among the antitriplet–antitriplet heavy baryons with vector mesons.

$$\begin{aligned}
\Pi^{\Xi_b^{\prime 0} \rightarrow \Sigma_b^0 \bar{K}^{*0}} &= \Pi_1^{(1)}(d, u, b) , \\
\Pi^{\Xi_b^{\prime -} \rightarrow \Sigma_b^- \bar{K}^{*0}} &= \sqrt{2} \Pi_1^{(1)}(d, d, b) , \\
\Pi^{\Omega_b^- \rightarrow \Xi_b^- \bar{K}^{*0}} &= \sqrt{2} \Pi_1^{(1)}(s, s, b) , \\
\Pi^{\Sigma_b^0 \rightarrow \Xi_b^0 K^{*0}} &= \Pi_1^{(1)}(d, u, b) , \\
\Pi^{\Sigma_b^- \rightarrow \Xi_b^- K^{*0}} &= \sqrt{2} \Pi_1^{(1)}(d, d, b) , \\
\Pi^{\Xi_b^{\prime -} \rightarrow \Omega_b^- K^{*0}} &= \sqrt{2} \Pi_1^{(1)}(s, s, b) , \\
\Pi^{\Sigma_b^0 \rightarrow \Sigma_b^0 \omega} &= \frac{1}{\sqrt{2}} \left[ \Pi_1^{(1)}(u, d, b) + \Pi_1^{(1)}(d, u, b) \right] , \\
\Pi^{\Sigma_b^+ \rightarrow \Sigma_b^+ \omega} &= \sqrt{2} \Pi_1^{(1)}(u, u, b) , \\
\Pi^{\Sigma_b^- \rightarrow \Sigma_b^- \omega} &= \sqrt{2} \Pi_1^{(1)}(d, d, b) ,
\end{aligned}$$

$$\begin{aligned}
\Pi^{\Xi_b^{\prime 0} \rightarrow \Xi_b^{\prime 0} \omega} &= \frac{1}{\sqrt{2}} \Pi_1^{(1)}(u, s, b) , \\
\Pi^{\Xi_b^{\prime -} \rightarrow \Xi_b^{\prime -} \omega} &= \frac{1}{\sqrt{2}} \Pi_1^{(1)}(d, s, b) , \\
\Pi^{\Xi_b^{\prime 0} \rightarrow \Xi_b^{\prime 0} \phi} &= \Pi_1^{(1)}(s, u, b) , \\
\Pi^{\Xi_b^{\prime -} \rightarrow \Xi_b^{\prime -} \phi} &= \Pi_1^{(1)}(s, d, b) , \\
\Pi^{\Omega_b^- \rightarrow \Omega_b^- \phi} &= 2\Pi_1^{(1)}(s, s, b) .
\end{aligned} \tag{A.1}$$

- Correlation functions responsible for the sextet–antitriplet transitions.

$$\begin{aligned}
\Pi^{\Xi_b^{\prime 0} \rightarrow \Xi_b^0 \rho^0} &= \frac{1}{\sqrt{2}} \Pi_2^{(1)}(u, s, b) , \\
\Pi^{\Xi_b^{\prime -} \rightarrow \Xi_b^- \rho^0} &= -\frac{1}{\sqrt{2}} \Pi_2^{(1)}(d, s, b) , \\
\Pi^{\Sigma_b^0 \rightarrow \Lambda_b \rho^0} &= \frac{1}{\sqrt{2}} \left[ \Pi_2^{(1)}(u, d, b) - \Pi_2^{(1)}(d, u, b) \right] , \\
\Pi^{\Sigma_b^- \rightarrow \Lambda_b \rho^-} &= \sqrt{2} \Pi_2^{(1)}(u, d, b) , \\
\Pi^{\Xi_b^{\prime -} \rightarrow \Xi_b^- \rho^-} &= \Pi_2^{(1)}(d, s, b) , \\
\Pi^{\Sigma_b^+ \rightarrow \Lambda_b \rho^+} &= -\sqrt{2} \Pi_2^{(1)}(d, u, b) , \\
\Pi^{\Xi_b^{\prime 0} \rightarrow \Xi_b^- \rho^+} &= \Pi_2^{(1)}(u, s, b) , \\
\Pi^{\Sigma_b^0 \rightarrow \Xi_b^0 \bar{K}^{*0}} &= -\Pi_2^{(1)}(d, u, b) , \\
\Pi^{\Sigma_b^- \rightarrow \Xi_b^- \bar{K}^{*0}} &= -\sqrt{2} \Pi_2^{(1)}(d, d, b) , \\
\Pi^{\Omega_b^- \rightarrow \Xi_b^- \bar{K}^{*0}} &= \sqrt{2} \Pi_2^{(1)}(s, s, b) , \\
\Pi^{\Xi_b^{\prime 0} \rightarrow \Lambda_b \bar{K}^{*0}} &= -\Pi_2^{(1)}(d, u, b) , \\
\Pi^{\Sigma_b^0 \rightarrow \Xi_b^0 K^{*0}} &= -\Pi_2^{(1)}(d, u, b) , \\
\Pi^{\Sigma_b^- \rightarrow \Xi_b^- K^{*0}} &= -\sqrt{2} \Pi_2^{(1)}(d, d, b) , \\
\Pi^{\Omega_b^- \rightarrow \Xi_b^- K^{*0}} &= \sqrt{2} \Pi_2^{(1)}(s, s, b) , \\
\Pi^{\Xi_b^{\prime 0} \rightarrow \Lambda_b K^{*0}} &= -\Pi_2^{(1)}(d, u, b) , \\
\Pi^{\Sigma_b^+ \rightarrow \Lambda_b K^{*+}} &= -\sqrt{2} \Pi_2^{(1)}(u, u, b) , \\
\Pi^{\Sigma_b^0 \rightarrow \Xi_b^- K^{*+}} &= -\Pi_2^{(1)}(u, d, b) , \\
\Pi^{\Xi_b^{\prime 0} \rightarrow \Xi_b^- K^{*+}} &= \Pi_2^{(1)}(d, s, b) , \\
\Pi^{\Sigma_b^- \rightarrow \Lambda_b K^{*-}} &= \sqrt{2} \Pi_2^{(1)}(d, d, b) , \\
\Pi^{\Omega_b^- \rightarrow \Xi_b^0 K^{*-}} &= \sqrt{2} \Pi_2^{(1)}(s, s, b) , \\
\Pi^{\Xi_b^{\prime -} \rightarrow \Xi_b^0 K^{*-}} &= \Pi_2^{(1)}(u, s, b) , \\
\Pi^{\Xi_b^{\prime 0} \rightarrow \Xi_b^0 \omega} &= \frac{1}{\sqrt{2}} \Pi_2^{(1)}(d, s, b) ,
\end{aligned}$$

$$\begin{aligned}
\Pi^{\Xi_b'^- \rightarrow \Xi_b^- \omega} &= \frac{1}{\sqrt{2}} \Pi_2^{(1)}(d, s, b) , \\
\Pi^{\Sigma_b^0 \rightarrow \Lambda_b \omega} &= \frac{1}{\sqrt{2}} \left[ \Pi_2^{(1)}(u, d, b) - \Pi_2^{(1)}(d, u, b) \right] , \\
\Pi^{\Xi_b'^0 \rightarrow \Xi_b^0 \phi} &= -\Pi_2^{(1)}(s, u, b) , \\
\Pi^{\Xi_b'^- \rightarrow \Xi_b^- \phi} &= -\Pi_2^{(1)}(s, d, b) .
\end{aligned} \tag{A.2}$$

- Correlation functions responsible for the antitriplet–antitriplet transitions.

$$\begin{aligned}
\Pi^{\Xi_b^0 \rightarrow \Xi_b^0 \rho^0} &= \frac{1}{\sqrt{2}} \Pi_3^{(1)}(u, s, b) , \\
\Pi^{\Xi_b^- \rightarrow \Xi_b^- \rho^0} &= -\frac{1}{\sqrt{2}} \Pi_3^{(1)}(d, s, b) , \\
\Pi^{\Lambda_b \rightarrow \Lambda_b \rho^0} &= -\frac{1}{\sqrt{2}} \left[ \Pi_3^{(1)}(d, u, b) - \Pi_3^{(1)}(u, d, b) \right] , \\
\Pi^{\Xi_b^- \rightarrow \Xi_b^0 \rho^-} &= \Pi_3^{(1)}(d, s, b) , \\
\Pi^{\Xi_b^0 \rightarrow \Xi_b^- \rho^+} &= \Pi_3^{(1)}(u, s, b) , \\
\Pi^{\Xi_b^0 \rightarrow \Lambda_b \bar{K}^{*0}} &= \Pi_3^{(1)}(u, u, b) , \\
\Pi^{\Xi_b^0 \rightarrow \Lambda_b K^{*0}} &= \Pi_3^{(1)}(u, u, b) , \\
\Pi^{\Xi_b^- \rightarrow \Lambda_b K^{*-}} &= -\Pi_3^{(1)}(u, d, b) , \\
\Pi^{\Xi_b^0 \rightarrow \Xi_b^0 \omega} &= \frac{1}{\sqrt{2}} \Pi_3^{(1)}(u, s, b) , \\
\Pi^{\Xi_b^- \rightarrow \Xi_b^- \omega} &= \frac{1}{\sqrt{2}} \Pi_3^{(1)}(d, s, b) , \\
\Pi^{\Lambda_b \rightarrow \Lambda_b \omega} &= \frac{1}{\sqrt{2}} \left[ \Pi_3^{(1)}(d, u, b) + \Pi_3^{(1)}(u, d, b) \right] , \\
\Pi^{\Xi_b^0 \rightarrow \Xi_b^0 \phi} &= \Pi_3^{(1)}(s, u, b) , \\
\Pi^{\Xi_b^- \rightarrow \Xi_b^- \phi} &= \Pi_3^{(1)}(s, d, b) .
\end{aligned} \tag{A.3}$$

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

## References

- [1] P. Biassoni, arXiv: 1009.2627 (2010).
- [2] G. Kane, (ed.), and A. Pierce, (ed.), “Perspectives on LHC physics”, (Michigan U.). 2008. 337pp. Hackensack, USA: World Scientific (2008) 337 p.
- [3] M. A. Shifman, A. I. Vainshtein and V. I. Zakharov, Nucl. Phys. B **147**, 385 (1979).
- [4] V. M. Braun, prep: hep-ph/9801222 (1998).
- [5] T. M. Aliev, K. Azizi, and M. Savcı, Phys. Rev. D **80**, 096003 (2009).
- [6] T. M. Aliev, A. Özpineci, S. B. Yakovlev, V. Zamiralov, Phys. Rev. D **74**, 116001 (2006).
- [7] T. M. Aliev, A. Özpineci, M. Savcı and V. Zamiralov, Phys. Rev. D **80**, 016010 (2009).
- [8] T. M. Aliev, K. Azizi, A. Özpineci and M. Savcı, Phys. Rev. D **80**, 096003 (2009).
- [9] T. M. Aliev, A. Özpineci, M. Savcı and V. Zamiralov, Phys. Rev. D **81**, 056004 (2010).
- [10] T. M. Aliev, K. Azizi, and M. Savcı, Nucl., Phys A, 847 (2010) 101.
- [11] P. Ball, V. M. Braun, Y. Koike, and K. Tanaka, Nucl. Phys. B **529**, 323 (1998).
- [12] P. Ball, V. M. Braun, Nucl. Phys. B **543**, 201 (1999); P. BALL, V. M. Braun, and A. Lenz, JHEP B **90**, 0708 (2007).
- [13] I. I. Balitsky and V. M. Braun, Nucl. Phys. B **311**, 239 (1988).
- [14] I. I. Balitsky, V. M. Braun, and A. V. Kolesnichenko, Nucl. Phys. B **312**, 509 (1989); K. G. Chetyrkin, A. Khodjamirian, and A. A. Pivovarov, Phys. Lett. B **651**, 250 (2008).
- [15] V. M. Belyaev, V. M. Braun, A. Khodjamirian and R. Rückl, Phys. Rev. D **51**, 6177 (1995).
- [16] T. M. Aliev, K. Azizi, A. Özpineci, Phys. Rev. D **79**, 056005 (2009).
- [17] P. Z. Huang, H. X. Chen, S. L. Zhu, Phys. Rev. D **80**, 094007 (2009).

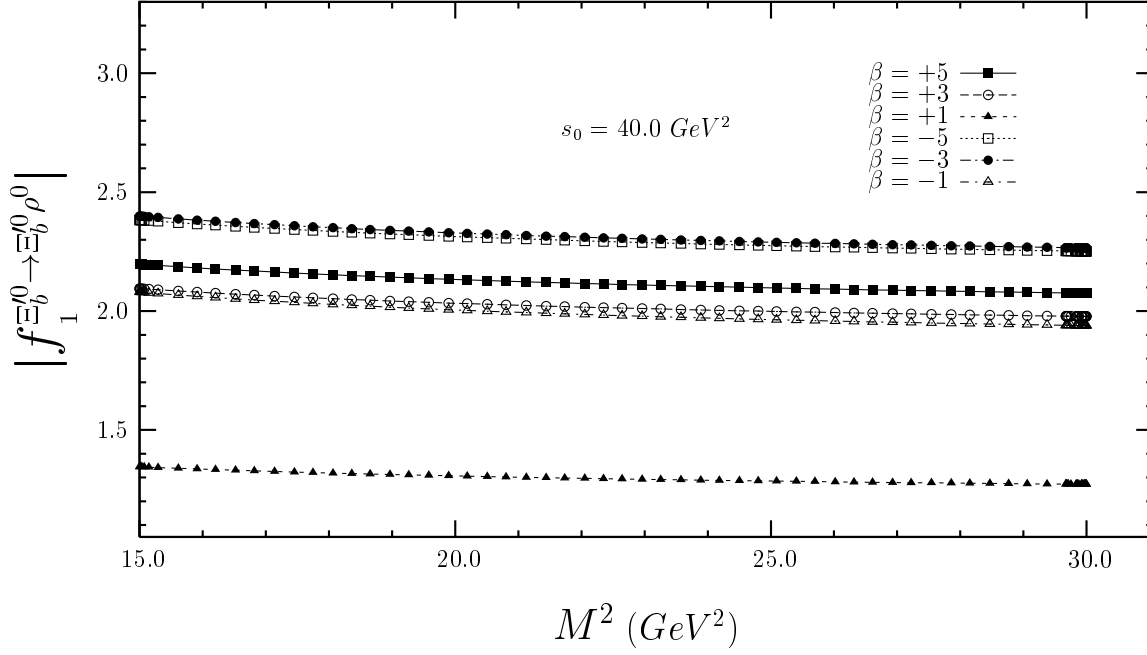


Figure 1: The dependence of the  $f_1^{\Xi_b'^0 \rightarrow \Xi_b'^0 \rho^0}$  on  $M^2$  at fixed values of the  $s_0$  and  $\beta$ .

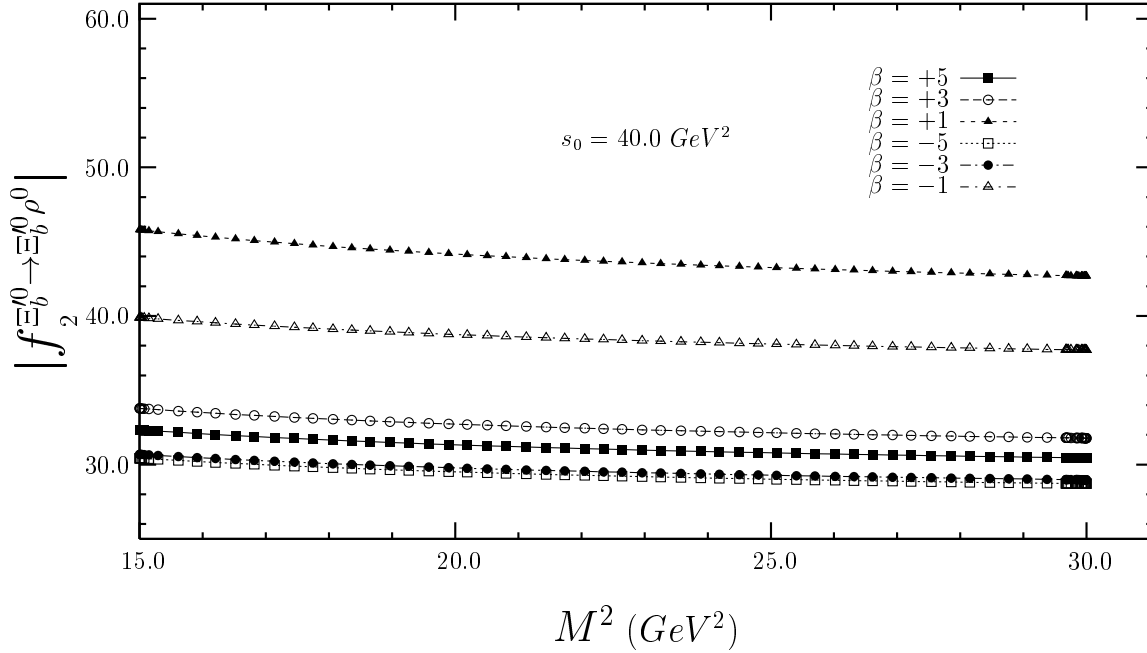


Figure 2: The dependence of the  $f_2^{\Xi_b'^0 \rightarrow \Xi_b'^0 \rho^0}$  on  $M^2$  at fixed values of the  $s_0$  and  $\beta$ .

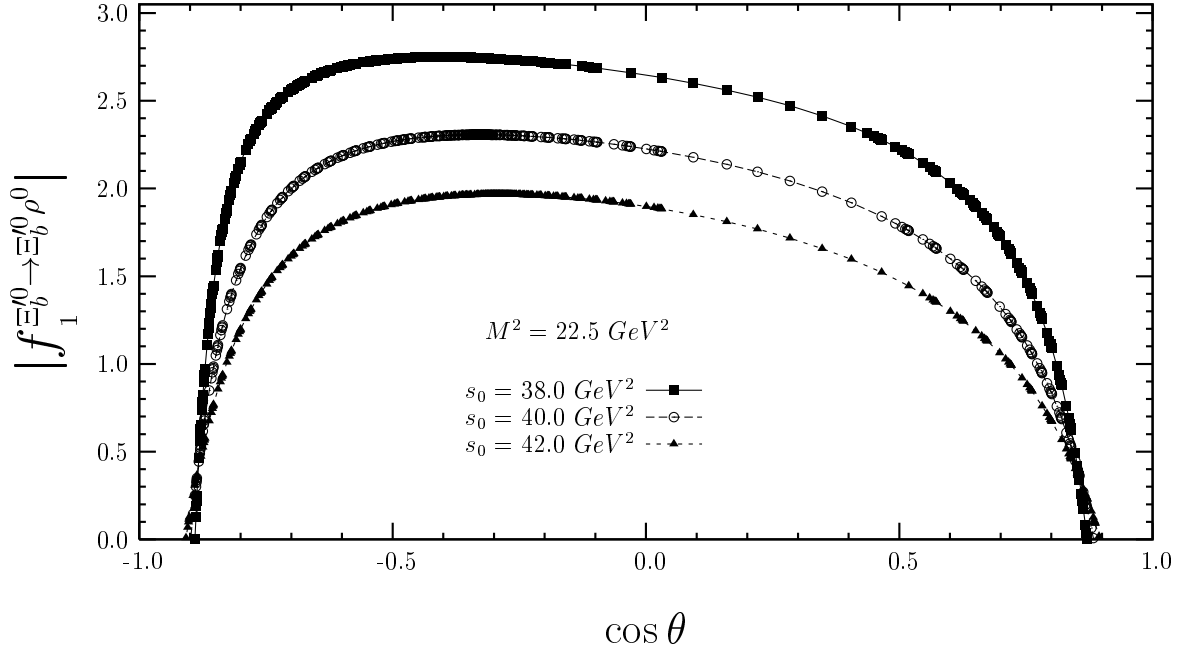


Figure 3: The dependence of the  $f_1^{\Xi_b'^0 \rightarrow \Xi_b'^0 \rho^0}$  on  $\cos\theta$  at fixed values of the  $s_0$  and  $M^2$ .

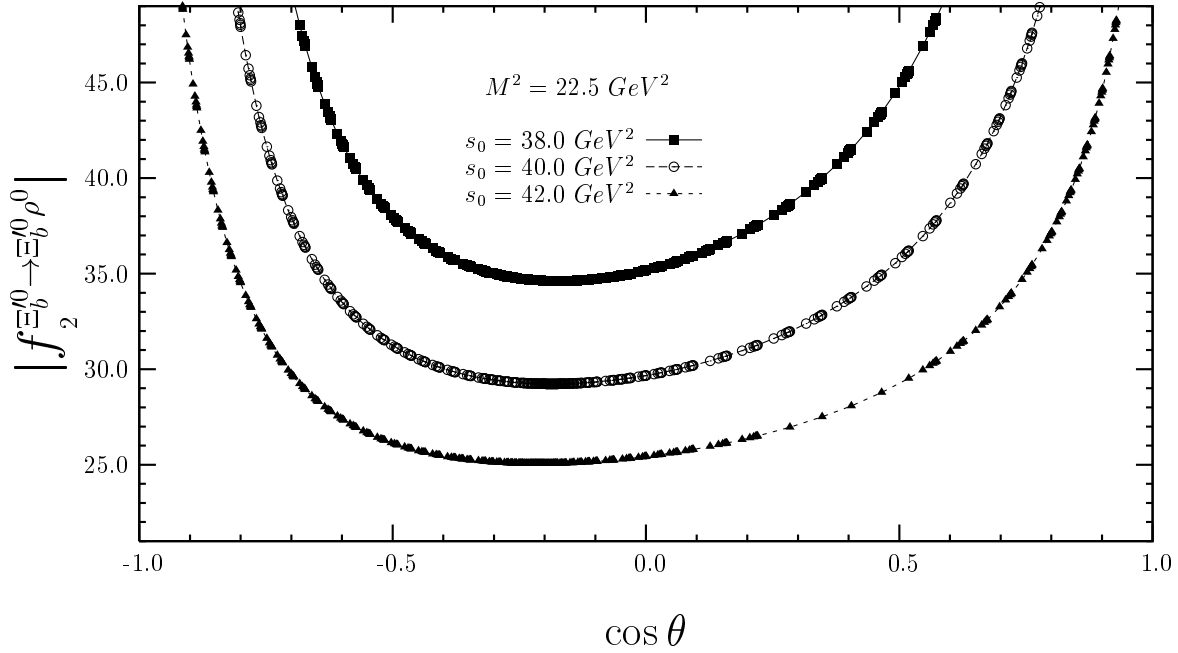


Figure 4: The dependence of the  $f_2^{\Xi_b'^0 \rightarrow \Xi_b'^0 \rho^0}$  on  $\cos\theta$  at fixed values of the  $s_0$  and  $M^2$ .