PHYSICAL REVIEW D **93,** 056007 (2016)

Analysis of the radiative decays $\Sigma_Q \to \Lambda_Q\gamma$ and $\Xi'_Q \to \Xi_Q\gamma$ in light cone sum rules

T. M. Aliev,[1,*,†] T. Barakat,[2,‡] and M. Savcı[1,§]

[1]Physics Department, Middle East Technical University, 06800 Ankara, Turkey
[2]Physics and Astronomy Department, King Saud University, P.O. Box 2454, Riyadh, Saudi Arabia

(Received 6 January 2016; published 18 March 2016)

The light cone sum rules method is used in studying the radiative decays $\Sigma_Q \to \Lambda_Q\gamma$ and $\Xi'_Q \to \Xi_Q\gamma$. First, the sum rules for the form factor $F_2(Q^2 = 0)$ responsible for these transitions is constructed. Using this result, the decay widths of the above-mentioned decays are calculated and analyzed. A comparison of our predictions on the decay widths of considered transitions with the predictions of the other approaches is presented.

DOI: 10.1103/PhysRevD.93.056007

I. INTRODUCTION

In the past decade, revolutionary progress has been made in hadron spectroscopy. Many baryons with single heavy quarks have been observed experimentally [1]. At the same time, many new charmoniumlike and bottomoniumlike states have also been discovered [2]. These states have more complicated structures compared to the ones predicted by the quark model.

These experimental achievements have carried the studies to a new level, namely, the study of the decays of these baryons. In this work, we concentrate our attention on the heavy baryons with single heavy quarks and investigate their electromagnetic decays.

According to $SU(3)$ classification, the heavy baryon ground states with spin $1/2$ belong to the sextet representation, spin $3/2$ to the sextet representation, and spin $1/2$ to the antitriplet representation. It is customary to denote these representations as 6, 6^*, and $\bar{3}$.

The radiative decays among the baryons belonging to these representations have already been studied in the framework of the nonrelativistic potential model [3] and light cone QCD sum rules incorporating the heavy quark effective theory [4], and incorporating both with heavy and chiral symmetry [5], $(2+1)$ flavor lattice QCD [6], heavy hadron chiral perturbation theory [7,8], chiral perturbation theory [9], the relativistic three-quark model [10], heavy quark symmetry [11], the static quark model [12], and the bag model [13], which lead to quite different results. Therefore, further independent calculations on these decay widths are necessary.

The present work is devoted to the study of the $\Sigma_Q \to \Lambda_Q\gamma$ and $\Xi'_Q \to \Xi_Q\gamma$ decays in the framework of the light cone QCD sum rules. Note that the decay widths between the $6^* \to 6$ and $6^* \to \bar{3}$ transitions have been studied earlier in the framework of the light cone QCD sum rules method in [14,15].

The paper is organized as follows. In Sec. II, the light cone QCD sum rules for the electromagnetic form factor $F_2(q^2 = 0)$ responsible for these decays are derived. Section III is devoted to the numerical analysis. In this section, we also present a comparison of our predictions with the results of other approaches.

II. LIGHT CONE QCD SUM RULES FOR THE $\Sigma_Q \to \Lambda_Q\gamma$ AND $\Xi'_Q \to \Xi_Q\gamma$ DECAY FORM FACTORS

In this section, we derive the light cone QCD sum rules for the radiative $\Sigma_Q \to \Lambda_Q\gamma$ and $\Xi'_Q \to \Xi_Q\gamma$ decay form factors. For this purpose, we start with the definition of the transition matrix element between heavy baryon states in the presence of the electromagnetic field, i.e., $\langle B_{Q_2}(p,s')|j_\mu^{el}|B_{Q_1}(p+q,s)\rangle$. This matrix element is parametrized in terms of the Dirac $F_1(Q^2)$ and Pauli $F_2(Q^2)$ form factors as follows:

$$
\begin{aligned}
&\langle B_{Q_2}(p,s')|j_\mu^{el}|B_{Q_1}(p+q,s)\rangle \\
&\quad = \bar{u}(p)\left[\left(\gamma_\mu - \frac{\not{q} q_\mu}{q^2}\right)F_1(Q^2) \right. \\
&\qquad \left. - \frac{1}{m_{B_{Q_1}} + m_{B_{Q_2}}} i\sigma_{\mu\nu}q^\nu F_2(Q^2)\right]u(p+q). \qquad (1)
\end{aligned}
$$

For the real photons, obviously, we need to know the values of these form factors only at the point $Q^2 = -q^2 = 0$. This process, i.e., transition of one of the heavy baryons in the sextet representation to another heavy baryon in the antitriplet representation in the presence of the electromagnetic field, is described by the following correlation function:

*taliev@metu.edu.tr
†Permanent address: Institute of Physics, Baku, Azerbaijan.
‡tbarakat@KSU.EDU.SA
§savci@metu.edu.tr

$$\Pi_\mu(p,q) = -\int d^4x \int d^4y e^{i(px+qy)} \times \langle 0|T\{\eta^a_{Q_1}(0) j_\mu(y) \bar{\eta}^s_{Q_2}(x)\}|0\rangle, \quad (2)$$

where $j_\mu = e_q \bar{q}\gamma_\mu q + e_Q \bar{Q}\gamma_\mu Q$ is the electromagnetic current with the electric charges e_q and e_Q for the light and heavy quarks, respectively, and $\eta^s_{Q_2}$ and $\eta^a_{Q_2}$ are the interpolating currents in the sextet and antitriplet representations, respectively.

The general form of the interpolating currents of the spin-1/2 heavy baryons in the sextet and antitriplet representations are given as (see for example [16])

$$\begin{aligned}
\eta^s_Q &= -\frac{1}{\sqrt{2}}\epsilon^{abc}\{(q_1^{aT} C Q^b)\gamma_5 q_2^c - (Q^{aT} C q_2^b)\gamma_5 q_1^c \\
&\quad + \beta(q_1^{aT} C\gamma_5 Q^b) q_2^c - \beta(Q^{aT} C\gamma_5 q_2^b) q_1^c\}, \\
\eta^a_Q &= \frac{1}{\sqrt{6}}\epsilon^{abc}\{2(q_1^{aT} C q_2^b)\gamma_5 Q^c + (q_1^{aT} C Q^b)\gamma_5 q_2^c \\
&\quad + (Q^{aT} C q_2^b)\gamma_5 q_1^c + 2\beta(q_1^{aT} C\gamma_5 q_2^b) Q^c \\
&\quad + (q_1^{aT} C\gamma_5 Q^b) q_2^c + (Q^{aT} C\gamma_5 q_2^b) q_1^c\}, \quad (3)
\end{aligned}$$

where β is the arbitrary auxiliary parameter, and the light quark contents of the heavy baryons in sextet and antitriplet representations are summarized in Table 1.

Introducing a plane wave electromagnetic background field $F_{\mu\nu} = i(\varepsilon_\mu q_\nu - \varepsilon_\nu q_\mu)e^{iqx}$, it is possible to rewrite the correlator (2) as follows:

$$\Pi_\mu(p,q)\varepsilon^\mu = i\int d^4x e^{ipx}\langle 0|T\{\eta_{Q_1}(0)\bar{\eta}_{Q_2}(x)\}|0\rangle_F, \quad (4)$$

where the subscript F means that all vacuum expectation values should be evaluated in the background field $F_{\mu\nu}$. The correlation function given in Eq. (2) can be obtained from Eq. (3) by expanding it in powers of the background field and considering only the linear term in $F_{\mu\nu}$, which corresponds to the single photon emission. More about the details of the application of the background field method can be found in [17] and [18].

In order to obtain the sum rules for the electromagnetic form factors describing the $\Sigma_Q \to \Lambda_Q\gamma$ and $\Xi'_Q \to \Xi_Q\gamma$ transitions, the correlation function is calculated in terms of hadrons from one side and in terms of the quark-gluon degrees of freedom by using the operator product expansion (OPE) and introducing the photon distribution amplitudes (DAs) from the other side. The photon DAs are the main nonperturbative ingredient of the light cone sum rules. In this version of the light cone QCD sum rules, OPE is performed by a twist of the nonlocal operator rather than dimensions of the operators in the traditional sum rules. The sum rules are obtained by matching these two representations.

TABLE I. Light quark contents of the heavy baryons in the symmetric sextet and antisymmetric antitriplet representations.

	$\Sigma^{+(++)}_{b(c)}$	$\Sigma^{0(+)}_{b(c)}$	$\Sigma^{-(0)}_{b(c)}$	$\Xi'^{-(0)}_{b(c)}$	$\Xi'^{0(+)}_{b(c)}$	$\Lambda^{0(+)}_{b(c)}$	$\Xi^{-(0)}_{b(c)}$	$\Xi^{0(+)}_{b(c)}$
q_1	u	u	d	d	u	u	d	u
q_2	u	d	d	s	s	d	s	s

We start our analysis by constructing the correlation function from the hadronic side. It can be obtained by inserting all intermediate hadronic sum rules, having the same quantum numbers as the corresponding interpolating currents η_Q. After isolating the ground state's contribution, we get

$$\Pi_\mu(p,q) = \frac{\langle 0|\eta^a_{Q_2}|B_{Q_2}(p_2)\rangle}{p_2^2 - m^2_{B_{Q_2}}}\langle B_{Q_2}(p_2)|j^{el}_\mu(q)|B_{Q_1}(p_1)\rangle \times \frac{\langle B_{Q_1}(p_1)|\bar{\eta}^s_{Q_1}|0\rangle}{p_1^2 - m^2_{B_{Q_1}}} + \cdots, \quad (5)$$

where the contributions coming from the higher states and continuum are denoted by dots, and $p_1 = p_2 + q$.

The expression for the correlator function can be obtained from the hadronic side by substituting the matrix elements appearing in Eq. (2). These matrix elements are defined in the standard way as follows:

$$\begin{aligned}
\langle 0|\eta_{B_{Q_2}}|B_{Q_2}(p_2)\rangle &= \lambda_2 u_{B_{Q_2}}(p_2), \\
\langle B_{Q_1}(p_1)|\eta_{B_{Q_1}}|0\rangle &= \lambda_1 \bar{u}_{B_{Q_1}}(p_1), \\
\langle B_{Q_2}(p_2)|j^{el}_\mu(q)|B_{Q_1}(p_1)\rangle &= \bar{u}_{B_{Q_2}}(p_2)\left[\left(\gamma_\mu - \frac{\not{q} q_\mu}{q^2}\right)F_1 - \frac{i\sigma_{\mu\nu}q^\nu}{m_{B_{Q_1}} + m_{B_{Q_1}}}F_2\right]u_{B_{Q_1}}(p_1),
\end{aligned} \quad (6)$$

where λ_i are the residues of the hadrons, B_{Q_i} are the baryons and $m_{B_{Q_i}}$ are their respective masses, and F_1 and F_2 are the Dirac and Pauli form factors, respectively. Using the equation of motion, the matrix element $\langle B_{Q_2}(p_2)|j^{el}_\mu(q)|B_{Q_1}(p_1)\rangle$ can be written as follows:

$$\begin{aligned}
&\langle B_{Q_2}(p_2)|j^{el}_\mu(q)|B_{Q_1}(p_1)\rangle \\
&= \bar{u}_{B_{Q_2}}(p_2)\left[\gamma_\mu(F_1 + F_2) - \frac{\not{q} q_\mu}{q^2}F_1 - \frac{(p_1+p_2)_\mu}{m_{B_{Q_1}} + m_{B_{Q_1}}}F_2\right]u_{B_{Q_1}}(p_1).
\end{aligned} \quad (7)$$

Inserting Eqs. (6) and (7) into Eq. (5) and performing summation over spins of the Dirac spinors, we get

$$\varepsilon^\mu \Pi_\mu(p,q) = \frac{\lambda_{B_{Q_1}} \lambda_{B_{Q_2}}}{(p^2 - m_{B_{Q_2}}^2)(p^2 - m_{B_{Q_1}}^2)} \times (\not{p} + m_{B_{Q_2}}) \left[\not{\varepsilon}(F_1 + F_2) - \frac{2(p\varepsilon)}{m_{B_{Q_1}} + m_{B_{Q_2}}} F_2 \right] (\not{p} + \not{q} + m_{B_{Q_1}}), \tag{8}$$

where we set $p_2 = p$, $p_1 = p + q$, and $q\varepsilon = 0$. It is easily seen from Eq. (8) that the correlation function possesses many structures, and any of them can be used for constructing the sum rules for the form factors $F_1 + F_2$ and F_2. The experience in working with the sum rules shows that the structures containing the maximum number of momenta exhibit rather good convergence. For this reason, in the calculation of the form factors $F_1 + F_2$ and F_2, we chose the structures $\not{p}\not{\varepsilon}\not{q}$ and $\not{p}(p\varepsilon)$, respectively. In this work, we calculate only the form factor F_2 since the transitions under consideration are described only by the form factor F_2. Note that the form factor $F_1 + F_2$ has already been calculated for the transitions under consideration in [19] and [20]. The expression of the correlator function given in Eq. (3) can be obtained in the deep Euclidean region in terms of photon DAs with increasing twist, where $p \ll 0$ and $(p+q)^2 \ll 0$.

Calculation of the correlation function can be carried out straightforwardly using the Wick's theorem. In performing this calculation, the expressions of the light and heavy quark propagators in the presence of the external field are needed. The light quark propagator in the background field is calculated in [21], and it is found that the contributions of the nonlocal operators $\bar{q}Gq$, $\bar{q}G^2q$, and $\bar{q}q\bar{q}q$ are quite small. Neglecting these contributions the expression of the light quark propagator can be written as

$$S_q(x) = \frac{i\not{x}}{2\pi^2 x^4} - \frac{m_q}{4\pi^2 x^2} - \frac{\langle \bar{q}q \rangle}{12}\left(1 - i\frac{m_q}{4}\not{x}\right) - \frac{x^2}{192} m_0^2 \langle \bar{q}q \rangle \left(1 - i\frac{m_q}{6}\not{x}\right) - ig_s \int_0^1 du \left[\frac{\not{x}}{16\pi^2 x^2} G_{\mu\nu}(ux)\sigma_{\mu\nu} - \frac{i}{4\pi^2 x^2} ux^\mu G_{\mu\nu}(ux)\gamma^\nu - i\frac{m_q}{32\pi^2} G_{\mu\nu}\sigma^{\mu\nu}\left(\ln\frac{-x^2\Lambda^2}{4} + 2\gamma_E\right)\right], \tag{9}$$

where γ_E is the Euler constant, and Λ is the cutoff energy separating the perturbative and nonperturbative regions, whose value is calculated in [22] to be $\Lambda = (0.5 \pm 0.1)$ GeV.

The expression of the heavy quark propagator in the background field in the x representation is given as

$$S_Q(x) = \frac{m_Q^2}{4\pi^2}\left\{ \frac{K_1(m_Q\sqrt{-x^2})}{\sqrt{-x^2}} + i\frac{\not{x}}{(\sqrt{-x^2})^2} K_2(m_Q\sqrt{-x^2}) \right\} - \frac{g_s}{16\pi^2}\int_0^1 du G_{\mu\nu}(ux)\left[(\sigma^{\mu\nu}\not{x} + \not{x}\sigma^{\mu\nu}) \times \frac{K_1(m_Q\sqrt{-x^2})}{\sqrt{-x^2}} + 2\sigma^{\mu\nu} K_0(m_Q\sqrt{-x^2}) \right], \tag{10}$$

where $K_i(m_Q\sqrt{-x^2})$ are the modified Bessel functions.

Having the expressions of the light and heavy quark propagators at hand, calculation of the theoretical part of the correlation function is a straightforward but rather tedious calculation. At this point, one technical remark is in order. To be able to express the vacuum expectation value $\langle 0|q(x)\bar{q}(0)|0\rangle_F$ in terms of the photon DAs, the Fierz identity needs to be used. It should be noted here that our approach in calculating the nonperturbative contribution to the correlation function follows the line of [23] for the $D^* D\pi$ coupling with the replacement of the pion DAs by the photon DAs.

As has already been noted in constructing the sum rules for the form factor $F_2(0)$, we have decided to choose the structure $(\varepsilon \cdot p)\not{p}\not{q}$ in both representations of the correlation function. In obtaining the final result for the sum rule of the form factor $F_2(0)$, the Borel transformation over the variables p^2 and $(p+q)^2$ is implemented using the quark-hadron duality ansatz. Using these steps of the calculation, we finally get the following sum rule for the form factor $F_2(0)$:

$$\lambda_{B_{Q_1}} \lambda_{B_{Q_2}} e^{-\left(\frac{m_{B_{Q_1}}^2}{M_1^2} + \frac{m_{B_{Q_2}}^2}{M_2^2}\right)} F_2(0) + \int ds_1 ds_2 \rho^h(s_1,s_2) e^{-\left(\frac{s_1}{M_1^2} + \frac{s_2}{M_2^2}\right)} = \Pi^{B(\mathrm{theor})}, \tag{11}$$

where $\lambda_{B_{Q_1}}$ and $\lambda_{B_{Q_2}}$ are the residues of the corresponding sextet and antitriplet baryons, respectively, whose expressions can be found in [19] and [20]; M_1^2 and M_2^2 are the Borel mass parameters for the corresponding channels. It should be noted here that, for consistency, the perturbative $\mathcal{O}(\alpha_s)$ corrections are neglected in the calculations of residues since they are not included in sum rules (11). These corrections might give considerable contribution to the form factor $F_2(0)$ similar to the $D^* D\pi$ case, but calculation of the radiative corrections lies beyond the scope of the present work. Explicit expression of $\Pi^{B(\mathrm{theor})}$ can be found in the Appendix. The second term on the left-hand side of Eq. (11) describes the contributions of the higher states and continuum. In calculating the contributions of these states, we use the quark-hadron duality ansatz; i.e., above some predetermined thresholds in the (s_1, s_2) plane, the hadronic spectral density is replaced by the QCD spectral density $\rho^{\mathrm{QCD}}(s_1, s_2)$. Using this ansatz, the continuum subtraction can be carried out by the procedure explained in [23]. Leaving aside the technical details, in the case $M_1^2 = M_2^2 = 2M^2$ and $u_0 = 1/2$, the subtraction procedure can be performed by using the following formula:

$$M^{2n} e^{-m_Q^2/M^2} \to \frac{1}{\Gamma(n)} \int_{m_Q^2}^{s_0} ds e^{-s/M^2} (s - m_Q^2)^{n-1}, \quad (n \geq 1).$$

We see from the expression of $\Pi^{B(\mathrm{theor})}$ that the leading twist term $\varphi_\gamma(u_0)$ is proportional to $m_b^4 M^4$, and higher twist terms are proportional to $m_b^4 M^2$ or $m_b^2 M^2$. Therefore, higher twist terms that are suppressed by inverse powers of M^2 with respect to the leading ones remain unaffected. Therefore, the continuum subtraction procedure is not performed for the higher twist terms (for more detail see [23]). It should be noted here that, in principle, single dispersion integrals originating in the subtractions, which make the double dispersion integral finite, can enter into the spectral density, but these terms are all eliminated by the double Borel transformations.

The masses of the initial and final heavy baryons are quite close to each other; hence, we can set $M_1^2 = M_2^2 = 2M^2$, which naturally leads to $u_0 = 1/2$. In our numerical analysis, we use these values of M^2 and u_0.

At the end of this section, we present the formula needed to calculate the decay rate of transitions under consideration, whose expression is as follows:

$$\Gamma(B_{Q_1} \to B_{Q_2}\gamma) = \frac{4\alpha|\vec{q}|^3}{(m_{B_{Q_1}} + m_{B_{Q_2}})^2} |F_2(0)|^2, \quad (12)$$

where

$$|\vec{q}| = \frac{(m_{B_{Q_1}}^2 - m_{B_{Q_2}}^2)}{2m_{B_{Q_1}}}$$

is the magnitude of the photon momentum.

III. NUMERICAL RESULTS

In this section, we perform the numerical analysis using the sum rules for the form factor $F_2(0)$. The input parameters of the values in this calculation are as follows: The quark condensate $\langle\bar{u}u\rangle(\mu = 1\ \mathrm{GeV}) = -(0.243)^3\ \mathrm{GeV}^3$, $\langle\bar{s}s\rangle|_{\mu=1\ \mathrm{GeV}} = 0.8\langle\bar{u}u\rangle|_{\mu=1\ \mathrm{GeV}}$, and $m_0^2 = (0.8 \pm 0.2)\ \mathrm{GeV}^2$, which is obtained from the analysis of the two-point sum rules for the light baryons [24,25] and B, B^* [26], and $f_{3\gamma} = -0.0039\ \mathrm{GeV}^2$ [18]; and magnetic susceptibility χ, which is calculated in [27–29], where we use $\chi(\mu = 1\ \mathrm{GeV}) = -2.85\ \mathrm{GeV}^{-2}$ in the present work.

The sum rules for the form factor $F_2(0)$ also contain three auxiliary parameters, namely, the Borel mass parameter M^2, the arbitrary parameter, and the continuum threshold s_0. Obviously, any physical quantity must be independent of the above-mentioned auxiliary parameters. Therefore, we should find the regions of these parameters for which the form factor $F_2(0)$ shows no sensitivity to their variation. The continuum threshold is related to the mass of the first excited state. The energy needed to excite the particle from the ground state to the first excited state is equal to $(\sqrt{s_0} - m)$, where m is the mass of the baryon in its ground state. Usually, $(\sqrt{s_0} - m)$ varies in the interval 0.3 GeV to 0.8 GeV. The experimental values of the heavy baryons are reproduced quite well if the continuum threshold varies in the following regions:

$$\sqrt{s_0} = \begin{cases} (3.1 \pm 0.1)\ \mathrm{GeV}, & \text{for } \Sigma_c \to \Lambda_c \\ (3.2 \pm 0.1)\ \mathrm{GeV}, & \text{for } \Xi_c' \to \Xi_c \\ (6.6 \pm 0.2)\ \mathrm{GeV}, & \text{for } \Sigma_b \to \Lambda_b \\ (6.7 \pm 0.2)\ \mathrm{GeV}, & \text{for } \Xi_b' \to \Xi_b. \end{cases} \quad (13)$$

The upper and lower bounds of the Borel mass parameter M^2 are determined by imposing the following two conditions:

(1) The contributions of the higher states and continuum should be less than the contributions of the ground state.

(2) Contributions of the higher twist terms should be less than the contributions of the leading twist terms.

As a result of these two conditions, the "working regions" of the Borel parameter for the transitions under consideration are determined to be

$$\begin{aligned} &2.0\ \mathrm{GeV}^2 \leq M^2 \leq 3.0\ \mathrm{GeV}^2, && \text{for } \Sigma_c \to \Lambda_c\gamma, \\ &2.2\ \mathrm{GeV}^2 \leq M^2 \leq 3.4\ \mathrm{GeV}^2, && \text{for } \Xi_c' \to \Xi_c\gamma, \\ &5.0\ \mathrm{GeV}^2 \leq M^2 \leq 7.0\ \mathrm{GeV}^2, && \text{for } \Sigma_b \to \Lambda_b\gamma, \\ &5.0\ \mathrm{GeV}^2 \leq M^2 \leq 7.5\ \mathrm{GeV}^2, && \text{for } \Xi_b' \to \Xi_b\gamma. \end{aligned} \quad (14)$$

In order to find the working region of the arbitrary parameter β for the transitions under consideration, we have studied the dependence of $F_2(0)$ on $\cos\theta$, where $\beta = \tan\theta$, at several fixed values of the continuum threshold s_0 and Borel parameters M^2 chosen from the working regions given in Eqs. (13) and (14), respectively. Our numerical analysis shows that in the domain $-0.7 \leq \cos\theta \leq -0.4$, which is common for all the considered radiative decays, the form factor $F_2(0)$ is practically independent of the arbitrary parameter β, and we finally obtain the following values for the form factor $F_2(0)$:

TABLE II. Decay widths of the $\Sigma_Q \to \Lambda_Q\gamma$ and $\Xi'_Q \to \Xi_Q\gamma$ transitions (in units of KeV).

	This work	[3]	[4]	[5]	[7]	[8]	[10]	[11]	[12]	[13]
$\Sigma_c^+ \to \Lambda_c^+\gamma$	50.0 ± 17.0	60.55	60	91.5	–	164	60.7 ± 1.5	87	120	46
$\Xi_c'^+ \to \Xi_c^+\gamma$	8.5 ± 2.5	–	–	19.7	–	54	12.7 ± 1.5	–	14	10
$\Xi_c'^0 \to \Xi_c^0\gamma$	0.27 ± 0.06	–	–	0.4	1.2 ± 0.7	0.02	0.17 ± 0.02	–	0.33	0.0015
$\Sigma_b^0 \to \Lambda_b^0\gamma$	152.0 ± 60.0	–	–	–	–	287.65	–	–	–	–
$\Xi_b'^0 \to \Xi_b\gamma$	47.0 ± 21.0	–	–	–	–	–	–	–	–	–
$\Xi_b'^- \to \Xi_b^-\gamma$	3.3 ± 1.3	–	–	–	3.11 ± 1.8	–	–	–	–	–

$$F_2(0) = \begin{cases} (3.0 \pm 0.5) & \text{for } \Sigma_c^+ \to \Lambda_c^+\gamma \\ (2.5 \pm 0.4) & \text{for } \Xi_c'^+ \to \Xi_c^+\gamma \\ (0.45 \pm 0.05) & \text{for } \Xi_c'^0 \to \Xi_c^0\gamma \\ (10.0 \pm 2.0) & \text{for } \Sigma_b^0 \to \Lambda_b^0\gamma \\ (9.0 \pm 2.0) & \text{for } \Xi_b'^0 \to \Xi_b^0\gamma \\ (2.4 \pm 0.5) & \text{for } \Xi_b'^- \to \Xi_b^-\gamma. \end{cases}$$

Note that exact $SU(3)$ U-spin flavor symmetry forbids the $\Xi_Q'^0 \to \Xi_Q^0\gamma$ decay. The nonzero value of $F_2(0)$ for this decay indicates the violation of the aforementioned symmetry.

A few words about the uncertainty in determination of the form factor $F_2(0)$ are in order. The radiative $\mathcal{O}(\alpha_s)$ corrections can of course bring their own uncertainty in the calculation of the form factor $F_2(0)$, which is not taken into account in the present work. We estimate that the uncertainties are only coming from the errors in the values of the input parameters entering into the sum rules.

Having calculated the values of the form factor $F_2(0)$, we can easily calculate the values of the considered decay widths by using Eq. (14), and the results are summarized in Table II. In this table, for completeness, we also present the predictions on the decay widths calculated in other approaches, such as the nonrelativistic quark model [3], the QCD sum rules method [4], heavy hadron chiral perturbation theory ([5,7], and [8]), the relativistic quark model [10], heavy quark symmetry that is implemented with the light quark symmetry [11], the naive static quark model [12], and the bag model [13].

From the comparison of our results with those existing in the literature, we see that our predictions are closer to the predictions of the relativistic quark model, and in particular, our result for the $\Xi_b'^- \to \Xi_b^-\gamma$ transition coincides with the result of [7]. We also observe that there appears to be a considerable difference among our results and the predictions of the other approaches on the decay widths of the considered transitions. Of course, only the experimental measurements of these decays can play the "judge" for choosing the right "theory."

In conclusion, we calculate the form factor $F_2(0)$ for the $\Sigma_Q \to \Lambda_Q\gamma$ and $\Xi'_Q \to \Xi_Q\gamma$ transitions within the light cone QCD sum rules method. The corresponding decay widths are estimated by using these values of the form factor $F_2(0)$. Comparison of our predictions on decay widths with the results of other approaches is presented.

APPENDIX

In this appendix, we present the explicit form of the correlation function $\Pi^{B(\mathrm{theor})}$ for the form factor $F_2(0)$, which is determined from the coefficient of the $(\varepsilon \cdot p)\not{p}\not{q}$ structure.

$$\begin{aligned}
\Pi^{B(\mathrm{theor})} &= \frac{\sqrt{3}}{128\pi^4}(1-\beta^2)(e_s - e_u)m_b^3 M^4(\mathcal{I}_2 - 2m_b^2\mathcal{I}_3 + m_b^4\mathcal{I}_4) \\
&+ \frac{1}{16\sqrt{3}\pi^2}(1-\beta)^2\chi m_b^4 M^4(e_s\langle\bar{s}s\rangle - e_u\langle\bar{u}u\rangle)(\mathcal{I}_3 - m_b^2\mathcal{I}_4)\varphi_\gamma(u_0) \\
&+ \frac{1}{1536\sqrt{3}\pi^4}(1-\beta)m_b M^2\{(1+\beta)(e_s - e_u)\langle g_s^2 G^2\rangle(3\mathcal{I}_2 - 4m_b^2\mathcal{I}_3) + 64\beta\pi^2 m_b^3(e_u\langle\bar{s}s\rangle - e_s\langle\bar{u}u\rangle)\mathcal{I}_3 \\
&- 24(1-\beta)m_b^3\pi^2(e_s\langle\bar{s}s\rangle - e_u\langle\bar{u}u\rangle)\mathbb{A}(u_0)\mathcal{I}_3 - 64e_b m_b\pi^2(\langle\bar{s}s\rangle - \langle\bar{u}u\rangle)(\mathcal{I}_2 - m_b^2\mathcal{I}_3)\} \\
&+ \frac{1}{64\sqrt{3}\pi^2}(1-\beta)m_b^2 M^2(e_s\langle\bar{s}s\rangle - e_u\langle\bar{u}u\rangle)\{[(5+\beta)\mathcal{I}_2 - 4(2+\beta)m_b^2\mathcal{I}_3][i_2(\mathcal{S},1) - i_2(\mathcal{T}_4,1)] \\
&+ [(1+5\beta)\mathcal{I}_2 - 4(1+2\beta)m_b^2\mathcal{I}_3][i_2(\tilde{\mathcal{S}},1) + i_2(\mathcal{T}_2,1)] - 8(2+\beta)m_b^2\mathcal{I}_3\tilde{j}_2(h_\gamma) \\
&+ 2(1-\beta)[(\mathcal{I}_2 - 2m_b^2\mathcal{I}_3)i_2(\mathcal{T}_1,1) + \mathcal{I}_2 i_2(\mathcal{T}_3,1)] \\
&- 2[(3+\beta)i_2(\mathcal{S},v) + (1+3\beta)i_2(\tilde{\mathcal{S}},v)](\mathcal{I}_2 - 2m_b^2\mathcal{I}_3)
\end{aligned}$$

$$
\begin{aligned}
&-4[\beta\mathcal{I}_2-(1+\beta)m_b^2\mathcal{I}_3]i_2(\mathcal{T}_2,v)-4(1-\beta)\mathcal{I}_2 i_2(\mathcal{T}_3,v)+4[\mathcal{I}_2-(1+\beta)m_b^2\mathcal{I}_3]i_2(\mathcal{T}_4,v\}\\
&+\frac{1}{32\sqrt{3}\pi^2}(1-\beta)m_b^3M^2f_{3\gamma}(e_s-e_u)[2(3+\beta)(\mathcal{I}_2-m_b^2\mathcal{I}_3)\tilde{j}_1(\psi^v)-(1+\beta)(\mathcal{I}_2-m_b^2\mathcal{I}_3)\psi^a(u_0)]\\
&+\frac{e^{-m_b^2/M^2}}{4608\sqrt{3}\pi^2M^2}(1-\beta)\langle g_s^2G^2\rangle(e_s\langle\bar{s}s\rangle-e_u\langle\bar{u}u\rangle)\{3(1+\beta)i_2(\mathcal{S},1)+3(1+\beta)i_2(\tilde{\mathcal{S}},1)\\
&+2i_2(\mathcal{T}_1,1)+3i_2(\mathcal{T}_2,1)-2i_2(\mathcal{T}_3,1)-3i_2(\mathcal{T}_4,1)-6i_2(\mathcal{S},v)-2i_2(\tilde{\mathcal{S}},v)-4i_2(\mathcal{T}_2,v)\\
&+4i_2(\mathcal{T}_3,v)+16\tilde{j}_2(h_\gamma)-\beta[2i_2(\mathcal{T}_1,1)-3i_2(\mathcal{T}_2,1)-2i_2(\mathcal{T}_3,1)+3i_2(\mathcal{T}_4,1)\\
&+2i_2(\mathcal{S},v)+6i_2(\tilde{\mathcal{S}},v)+4i_2(\mathcal{T}_3,v)-4i_2(\mathcal{T}_4,v)-8\tilde{j}_2(h_\gamma)]\}\\
&+\frac{e^{-m_b^2/M^2}}{4608\sqrt{3}\pi^2M^2}(1-\beta)\{(1-\beta)\langle g_s^2G^2\rangle(e_s\langle\bar{s}s\rangle-e_u\langle\bar{u}u\rangle)\mathbb{A}(u_0)\\
&-2(e_u\langle\bar{s}s\rangle-e_s\langle\bar{u}u\rangle)[3(1+\beta)\langle g_s^2G^2\rangle-(3+\beta)\langle g_s^2G^2\rangle\\
&-8(11+5\beta)f_{3\gamma}m_0^2\pi^2\tilde{j}_1(\psi^v)-4(2+5\beta)f_{3\gamma}m_0^2\pi^2\psi^a(u_0)]\}\\
&+\frac{e^{-m_b^2/M^2}}{96\sqrt{3}M^4}(1-\beta)f_{3\gamma}m_0^2m_b^2(e_u\langle\bar{s}s\rangle-e_s\langle\bar{u}u\rangle)[2(3+\beta)\tilde{j}_1(\psi^v)+\beta\psi^a(u_0)]\\
&+\frac{e^{-m_b^2/M^2}}{13824\sqrt{3}\pi^2M^6}(1-\beta)\langle g_s^2G^2\rangle m_b^2(e_u\langle\bar{s}s\rangle-e_s\langle\bar{u}u\rangle)\{3\beta m_0^2+8f_{3\gamma}\pi^2[2(3+\beta)\tilde{j}_1(\psi^v)+\beta\psi^a(u_0)]\}\\
&+\frac{e^{-m_b^2/M^2}}{3456\sqrt{3}M^8}(1-\beta)f_{3\gamma}\langle g_s^2G^2\rangle m_0^2m_b^2(e_u\langle\bar{s}s\rangle-e_s\langle\bar{u}u\rangle)[2(3+\beta)\tilde{j}_1(\psi^v)+\beta\psi^a(u_0)]\\
&-\frac{e^{-m_b^2/M^2}}{6912\sqrt{3}M^{10}}(1-\beta)f_{3\gamma}\langle g_s^2G^2\rangle m_0^2m_b^4(e_u\langle\bar{s}s\rangle-e_s\langle\bar{u}u\rangle)[2(3+\beta)\tilde{j}_1(\psi^v)+\beta\psi^a(u_0)]\\
&-\frac{e^{-m_b^2/M^2}}{64\sqrt{3}\pi^2}(1-\beta)\beta m_0^2(e_u\langle\bar{s}s\rangle-e_s\langle\bar{u}u\rangle)-\frac{1}{384\sqrt{3}\pi^2}(1-\beta)(2-\beta)m_0^2m_b^2[(e_u\langle\bar{s}s\rangle-e_s\langle\bar{u}u\rangle)+e_b(\langle\bar{s}s\rangle-\langle\bar{u}u\rangle)]\mathcal{I}_2\\
&-\frac{e^{-m_b^2/M^2}}{1152\sqrt{3}\pi^2m_b}(1-\beta)(3+\beta)f_{3\gamma}[(e_s-e_u)\langle g_s^2G^2\rangle\\
&+96m_b\pi^2(e_u\langle\bar{s}s\rangle-e_s\langle\bar{u}u\rangle)-3m_b^2e^{m_b^2/M^2}(e_s-e_u)\langle g_s^2G^2\rangle\mathcal{I}_2]\tilde{j}_1(\psi^v)\\
&-\frac{1}{1152\sqrt{3}\pi^2}(1-\beta)^2\langle g_s^2G^2\rangle m_b^2(e_s\langle\bar{s}s\rangle-e_u\langle\bar{u}u\rangle)\chi\mathcal{I}_2\varphi_\gamma(u_0)\\
&-\frac{e^{-m_b^2/M^2}}{2304\sqrt{3}}(1-\beta)f_{3\gamma}\{96\beta(e_u\langle\bar{s}s\rangle-e_s\langle\bar{u}u\rangle)-\frac{1}{\pi^2m_b}(1+\beta)(e_s-e_u)\langle g_s^2G^2\rangle(1-3m_b^2e^{m_b^2/M^2}\mathcal{I}_2)\}\psi^a(u_0).
\end{aligned}
$$

The functions $i_\ell(\phi,f(v))$, $\tilde{j}_\ell(f(u))$, where $(\ell=1,2)$, and $\mathcal{I}_n$ entering into the correlation function $\Pi^{B(\mathrm{theor})}$ are defined as

$$
\begin{aligned}
i_1(\phi,f(v))&=\int\mathcal{D}\alpha_i\int_0^1 dv\phi(\alpha_{\bar{q}},\alpha_q,\alpha_g)f(v)\delta'(k-u_0),\\
i_2(\phi,f(v))&=\int\mathcal{D}\alpha_i\int_0^1 dv\phi(\alpha_{\bar{q}},\alpha_q,\alpha_g)f(v)\delta''(k-u_0),\\
\tilde{j}_1(f(u))&=\int_{u_0}^1 du f(u),\\
\tilde{j}_2(f(u))&=\int_{u_0}^1 du(u-u_0)f(u),\\
\mathcal{I}_n&=\int_{m_b^2}^\infty ds\frac{e^{-s/M^2}}{s^n},
\end{aligned}
$$

where

$$k = \alpha_q + \alpha_g \bar{v}, \qquad u_0 = \frac{M_1^2}{M_1^2 + M_2^2}, \qquad M^2 = \frac{M_1^2 M_2^2}{M_1^2 + M_2^2}.$$

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