Semileptonic decay $\Lambda_c \rightarrow \Lambda \ell^+ \nu$ from QCD light-cone sum rules

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(Dated: March 26, 2022)

We present the study of the semileptonic decay $\Lambda_c \rightarrow \Lambda \ell^+ \nu$ by using the light-cone sum rule approach. Distribution amplitudes (DAs) for the $\Lambda$ baryon are discussed to the leading order conformal spin, and QCD sum rule estimate for the corresponding parameters is presented. The form factors describing the decay are calculated and used to predict the decay width and the decay asymmetry parameter $\alpha$. With the inclusion of twist-3 contributions the calculated decay width $\Gamma = (7.2 \pm 2.0) \times 10^{-14}$ GeV as well as asymmetry $\alpha = -(0.88 \pm 0.03)$ is found in good agreement with the experimental data, while there are appreciable deviations from experiment values when the higher twist contributions are included.

PACS numbers: 13.30.-a, 14.20.Lq, 11.55.Hx

I. INTRODUCTION

The study of flavor changing decays of the $c$-quark is always an active field in the heavy flavor physics, and the most obvious reason lies on that those processes can provide useful information on the various charm related Cabibbo-Kobayashi-Maskawa (CKM) matrix elements which are the main ingredients of the standard model (SM). Furthermore, the thorough understanding of the SM in itself needs a comprehension of the flavor changing dynamics. Unfortunately, such a comprehension is difficult contemporarily, the fact is that form factors characterizing those processes are not perturbative quantities and whose determination must invoke some non-perturbative method. This paper aims to give a preliminary determination of the form factors of semileptonic $\Lambda_c \rightarrow \Lambda \ell^+ \nu$ decay. In the calculation we will use the method of QCD sum rules on the light-cone $^{[1]}$, which in the past has been successfully applied to various problems in heavy meson physics, see $^{[2]}$ for a review.

The method of light-cone sum rules (LCSR) is a new development of the standard technique of QCD sum rules à la SVZ sum rules $^{[3]}$, which comes as the remedy for the conventional approach in which vacuum condensates carry no momentum $^{[4]}$. The main difference between SVZ sum rule and LCSR is that the short-distance Wilson OPE in increasing dimension is replaced by the light-cone expansion in terms of distribution amplitudes of increasing twist, originally used in the description of the hard exclusive process $^{[5]}$. In recent years there have been many applications of LCSR to baryons. The nucleon electromagnetic form factors were studied for the first time in $^{[6, 7]}$ and, more recently, in $^{[8, 9]}$ for a further consideration. Several nucleon related processes gave fruitful results within LCSR, the weak decay $\Lambda_b \rightarrow p\ell\nu_\ell$ was considered in $^{[10]}$ in both full QCD and HQET LCSR. The generalization to the $N\gamma\Delta$ transition form factor was worked out in $^{[11]}$.

In this paper we will adopt the LCSR approach to study the exclusive semileptonic decay $\Lambda_c \rightarrow \Lambda \ell^+ \nu$. This transition had been studied in the literature by several authors, employing flavor symmetry or quark model or both in Refs. $^{[12, 13, 14, 15, 16]}$. There are also QCD sum rule description of the form factors $^{[17]}$, upon which the total decay rate are obtained.

The paper is organized as follows: The relevant $\Lambda$ baryon DAs are first discussed in Sec. II. Following that Sec. III is devoted to the LCSRs for the semileptonic $\Lambda_c \rightarrow \Lambda \ell^+ \nu$ decay form factors. The numerical analysis and our conclusion are presented in Sec. IV.

II. THE $\Lambda$ BARYON DISTRIBUTION AMPLITUDES

Our discussion in this section for the $\Lambda$ baryon DAs parallels with that for the nucleon $^{[18]}$, so we only list the results following from that procedure and for details it is recommended to consult the original paper. The $\Lambda$ baryon
DAs are defined through the matrix element

\[ 4 \langle 0 | \epsilon_{ijk} u_i^\alpha(a_1 x) d_j^\beta(a_2 x) s_k^\gamma(a_3 x) | P \rangle = (A_1 + \frac{x^2 M^2}{4} A_1^M)(P \gamma_5 C)_{\alpha \beta} \Lambda_\gamma + A_2 M (P \gamma_5 C)_{\alpha \beta} (\not \Lambda) \gamma + A_3 M (\gamma^\mu \gamma_5 C)_{\alpha \beta} \Lambda_\gamma + A_4 M^2 (\gamma^\mu \gamma_5 C)_{\alpha \beta} (i \sigma^{\mu \nu} x_{\nu \Lambda}) \gamma + A_5 M^3 (\gamma^\mu \gamma_5 C)_{\alpha \beta} (\not \Lambda) \gamma, \]

where the \( \Lambda \) generically designates the spinor for the \( \Lambda \) baryon with momentum \( P \). Only axial-vector DAs are presented here, for those with other Lorentz structures do not contribute in the final sum rules. The twist of those calligraphic DAs is indefinite, but they can be related to the ones with definite twist as:

\[ A_1 = A_4, \quad 2 P \cdot x A_2 = -A_1 + A_2 - A_3, \]
\[ 2 A_3 = A_5, \quad 4 P \cdot x A_4 = -2 A_1 - A_3 - A_4 + 2 A_5, \]
\[ 4 P \cdot x A_5 = A_3 - A_4, \quad (2 P \cdot x)^2 A_6 = A_1 - A_2 + A_3 + A_4 - A_5 + A_6. \] (2)

The twist of \( A_i \) is given in Tab. I Each distribution amplitudes \( F = A_i \) can be represented as Fourier integral over the longitudinal momentum fractions \( x_1, x_2, x_3 \) carried by the quarks inside the baryon with \( \Sigma_i x_i = 1 \),

\[ F(a_i P \cdot x) = \int D x e^{-ip \cdot x \Sigma_i x_i a_i} F(x_i). \] (3)

The integration measure is defined as

\[ \int D x = \int_0^1 dx_1 dx_2 dx_3 \delta(x_1 + x_2 + x_3 - 1). \] (4)

As elucidated in [19], those distribution amplitudes are scale dependent and can be expanded into orthogonal functions with increasing conformal spin. To the leading conformal spin, or \( s \)-wave, accuracy the expansion reads [18]

\[
\begin{align*}
A_1(x_i, \mu) &= -120 x_1 x_2 x_3 \phi_0^0(\mu), \\
A_2(x_i, \mu) &= -24 x_1 x_2 \phi_1^0(\mu), \\
A_3(x_i, \mu) &= -12 x_3 (1 - x_3) \psi_0^0(\mu), \\
A_4(x_i, \mu) &= -3 (1 - x_3) \psi_0^1(\mu), \\
A_5(x_i, \mu) &= -6 x_3 \phi_1^1(\mu), \\
A_6(x_i, \mu) &= -2 \phi_0^0(\mu),
\end{align*}
\] (5)

where the constraint \( A(x_1, x_2, x_3) = A(x_2, x_1, x_3) \) arising from the condition that the \( \Lambda \) baryon has isospin 0 has been used in the derivation. All the 6 parameters involved in Eq. 5 can be expressed in terms of 2 independent matrix elements of local operators. Those parameters are

\[ \phi_0^0 = \phi_6^0 = -f_\Lambda, \quad \phi_1^0 = \phi_5^0 = \frac{1}{2} (\lambda_1 + f_\Lambda), \quad \psi_0^0 = \psi_5^0 = \frac{1}{2} (\lambda_1 - f_\Lambda). \]

The normalization of \( A_1 \) at the origin defines the nucleon coupling constant \( f_\Lambda \),

\[ \langle 0 | \epsilon_{ijk} u_i^\alpha(0) C \gamma_5 \not d_j(0) | P \rangle = f_\Lambda z \cdot P \not \Lambda(P). \] (6)

The remaining parameter \( \lambda_1 \) is defined by the matrix element

\[ \langle 0 | \epsilon_{ijk} u_i^\alpha(0) C \gamma_5 \gamma_\mu \not d_j(0) | P \rangle = \lambda_1 M \Lambda(P). \] (7)

The twist of the order \( O(x^2) \) correction starts from twist five, which is apparent in the definition 11. Due to the numerically small contribution it gives in the previous applications of light-cone QCD sum rules 10, we do not consider it in the following analysis.
Complying with the standard philosophy in the sum rule analysis, we consider the following correlation function

\[ z'^v T_\nu(P,q) = iz'^v \int d^4xe^{iux} \langle 0 | T\{ j_{\Lambda_c}(0)j_\nu(x) \} | P \rangle, \]  

where \( j_{\Lambda_c} = \epsilon_{ijk}(u^iC\gamma_5\ell d^j)\ell \gamma^k \) is the current interpolating the \( \Lambda_c \) baryon state, \( j_\nu = \bar{c}\gamma_\nu(1 - \gamma_5)s \) is the weak current, \( C \) is the charge conjugation matrix, and \( i, j, k \) denote the color indices. The auxiliary light-cone vector \( z \) is introduced to project out the main contribution on the light-cone. The interpolating current used here is not the unique one, as exemplified in case studies \( \text{[20]} \) for the applications of QCD sum rules, and there can be other choices. The coupling constant of the interpolating current to the vacuum can thus be defined as

\[ \langle 0 | j_{\Lambda_c} | \Lambda_c(P') \rangle = f_{\Lambda_c} z \cdot P' \Lambda_c(P'), \]

where \( \Lambda_c(P') \) and \( P' \) is the \( \Lambda_c \) baryon spinor and four-momentum, respectively. Form factors are given in the usual way

\[ \langle \Lambda_c(P - q) | j_\nu | \Lambda(P) \rangle = \bar{\Lambda}(P - q) \left[ f_1 \gamma_\nu - i \frac{f_2}{M_{\Lambda_c}} \sigma_{\nu\mu} q^\mu \right] \Lambda(P), \]

in which \( M_{\Lambda_c} \) is the \( \Lambda_c \) mass, \( \Lambda(P) \) denotes the \( \Lambda \) spinor, satisfying \( P\Lambda(P) = MA(P) \), where \( M \) is the \( \Lambda \) mass and \( P \) its four-momentum. Those form factors give no contribution in the case of massless final leptons are omitted here.

Giving those definitions, the hadronic representation of the correlation function \( \text{[8]} \) can be written as

\[ z'^v T_\nu = \frac{2f_{\Lambda_c}}{M_{\Lambda_c}^2 - P'^2} (z \cdot P')^2 \left[ f_1 \gamma_\nu + f_2 \frac{\gamma_\nu}{M_{\Lambda_c}} \right] \Lambda(P) + \cdots, \]

where \( P' = P - q \) and the dots stand for the higher resonances and continuum. While on the theoretical side, at large Euclidean momenta \( P'^2 \) and \( q^2 \) the correlation function \( \text{[8]} \) can be calculated perturbatively and the result, in the leading order of \( \alpha_s \), is

\[ z'^v T_\nu = -2(z \cdot P)^2 \int d^4x \left[ \frac{d^4k}{(2\pi)^4} \frac{z \cdot k}{k^2 - m_c^2} e^{i(k + q) \cdot x} \langle 0 | \epsilon_{ijk} u^i(0) d^j(0) s_\mu(x) | P \rangle \right], \]

where \( m_c \) is the c-quark mass. Substituting \( \text{[11]} \) into Eq. \( \text{[12]} \) we obtain,

\[ z'^v T_\nu = -2(z \cdot P)^2 \int d^4x \left[ \frac{x_3 B_0(x_3)}{M^2} + 2 M^4 \int d^4x \frac{x_3^2 B_2(x_3)}{k^2 - m_c^2} \right] \bar{\Lambda}(1 - \gamma_5) \Lambda(P) + \cdots, \]

where \( k = x_3 P - q \) and the ellipses stand for contributions that are nonleading in the infinite momentum frame kinematics \( P \to \infty, q \sim \text{const.}, z \sim 1/P \). The functions \( B_i \) are defined by

\[ B_0 = \int_0^{1-x_3} dx_1 x_1(1 - x_1 - x_3), \]

\[ B_1 = -2A_3 - A_4 - A_5, \]

\[ B_2 = A_1 - A_2 + A_3 + A_4 - A_5 + A_6, \]

\[ B_3 = -A_1 - A_2 - A_3. \]

The DAs with tildes are defined via integration as follows

\[ \tilde{A}(x_3) = \int_1^{x_3} dx_3' \int_0^{1-x_3'} dx_1 A(x_1, 1 - x_1, x_3'), \]

\[ \tilde{A}(x_3) = \int_1^{x_3} dx_3' \int_0^{x_3} dx_3'' \int_0^{1-x_3''} dx_1 A(x_1, 1 - x_1, x_3'', x_3'). \]
The origin of those functions can be traced back to the partial integration adopted to eliminate the factor $1/P \cdot x$ which appears in the distribution amplitudes. When the next-to-leading order conformal expansion is considered, the surface terms completely sum to zero. The term $B_0$ corresponds to the leading twist contribution. The form factors $f_2$ and $g_2$ in (13) are characterized by the higher twist contributions.

Equating (11) and (13), adopting the quark-hadron duality assumption and employing a Borel improvement of $P^2$ on both sides lead us to the desired sum rules for the form factors $f_1$ and $f_2$,

$$-f_{\lambda_1}f_1 e^{-M^2_{\lambda_1}/M_B^2} = - \int_{x_0}^1 dx_2 e^{-s'/M_B^2} \left[ B_0 + \frac{M^2}{M_B^2} \left( -B_1(x_3) + \frac{M^2}{M_B^2} B_2(x_3) \right) \right] + \frac{M^2 x_0^2 e^{-s_0/M_B^2}}{m^2 + Q^2 + x_0^2 M^2} \left[ B_1(x_0) - \frac{M^2}{M_B^2} x_0 B_2(x_0) \right] + \frac{M^2 e^{-s_0/M_B^2} x_0^2}{m^2 + Q^2 + x_0^2 M^2} \frac{d}{dx_0} \left( \frac{M^2 x_0^2 B_2(x_0)}{m^2 + Q^2 + x_0^2 M^2} \right),$$

and

$$\frac{f_{\lambda_1} f_2}{M_{\lambda_1} M} e^{-M^2_{\lambda_1}/M_B^2} = \frac{1}{M_B^2} \int_{x_0}^1 \frac{dx_3}{x_3} e^{-s'/M_B^2} \left( B_3(x_3) - \frac{M^2}{M_B^2} B_2(x_3) \right) + \frac{x_0 e^{-s_0/M_B^2}}{m^2 + Q^2 + x_0^2 M^2} \left( B_3(x_0) - \frac{M^2}{M_B^2} x_0 B_2(x_0) \right) + \frac{M^2 e^{-s_0/M_B^2} x_0^2}{m^2 + Q^2 + x_0^2 M^2} \frac{d}{dx_0} \left( \frac{x_0 B_2(x_0)}{m^2 + Q^2 + x_0^2 M^2} \right),$$

where

$$s' = (1-x)M^2 + \frac{m^2 + (1-x)Q^2}{x},$$

and $x_0$ is the positive solution of the quadratic equation for $s' = s_0$:

$$2M^2 x_0 = \sqrt{(Q^2 + s_0 - M^2)^2 + 4M^2(Q^2 + m^2) - (Q^2 + s_0 - M^2)}.$$

As the sum rules for the form factors $g_1$ and $g_2$ are identical with those for the $f_1$ and $f_2$, $f_1 = g_1$ and $f_2 = g_2$, we will only discuss the results for $f_1$ and $f_2$ in the following sections.

**IV. NUMERICAL ANALYSIS AND THE CONCLUSION**

**A. Values for $f_{\lambda_1}$ and $\lambda_1$**

Parameters $f_{\lambda_1}$, $\lambda_1$ appear in the conformal expansion of the DAs, so we have to determine their values before proceeding to analyze the LCSRs. According to their definitions, we consider correlation functions

$$\Pi(q^2) = i \int d^4xe^{iqx} \langle 0 | T \{ J_i(x), \bar{J}_j(0) \} | 0 \rangle.$$

where $J_i$'s are given in (9) and (17). Following the standard QCD sum rule philosophy, an estimate for $f_{\lambda_1}$ and their relative sign is straightforward

$$(4\pi)^4 f_{\lambda_1}^2 e^{-M^2_{\lambda_1}/M_B^2} = \frac{2}{5} \int_{m^2}^{s_0} s(1-x)^5 e^{-s/M_B^2} ds - \frac{b}{3} \int_{m^2}^{s_0} x(1-x)(1-2x)e^{-s/M_B^2} ds,$$

$$4(2\pi)^4 \lambda_1^2 M^2 e^{-M^2/M_B^2} = \frac{1}{2} \int_{m^2}^{s_0} s^2 \left[ (1-x^2)(1-8x + x^2) - 12x \ln x \right] e^{-s/M_B^2} ds$$

$$+ \frac{b}{12} \int_{m^2}^{s_0} (1-x)^2 e^{-s/M_B^2} ds - \frac{4}{3} a^2 e^{-M^2/M_B^2},$$

$$(2\pi)^4 f_{\lambda_1} \lambda_1 M e^{-M^2/M_B^2} = \frac{m_s}{6} \int_{m^2}^{s_0} s \left[ (1-x)(3 + 13x - 5x^2 + x^3) + 12x \ln x \right] e^{-s/M_B^2} ds$$

$$+ \frac{b}{12} \int_{m^2}^{s_0} (1-x) \left[ 1 + \frac{1}{3}(1-x)(5 - \frac{2}{x}) \right] e^{-s/M_B^2} ds.$$
where \( x = m_s^2/s \) and \( m_s \) is the s-quark mass. The sum rule for the \( f_\Lambda \) has been obtained before \[10\], where the corresponding heavy quark limit is also derived. At the working window \( s_0 \sim 1.6^2 \text{GeV}^2 \) and \( 1 < M_B^2 < 2 \text{GeV}^2 \) the numerical value for the coupling constant reads

\[
f_\Lambda = 6.1 \times 10^{-3} \text{GeV}^2, \quad \lambda_1 = -1.2 \times 10^{-2} \text{GeV}^2. \tag{24}
\]

The relative sign of \( \lambda_1/f_\Lambda \) is obtained from the sum rule \[23\]. In the numeric analysis, the standard values \( a = -(2\pi)^2\bar{q}q = 0.55 \text{GeV}^3 \), \( b = (2\pi)^2(\alpha_s G^2/\pi) = 0.47 \text{GeV}^4 \) and \( m_s = 0.15 \text{GeV} \) are adopted. It should be noted that our value for \( f_\Lambda \) here does coincide with that obtained in \[21\].

The sum rule for the coupling constant of \( \Lambda_c \) to vacuum is similar to that for \( f_\Lambda \), where the simple substitution \( m_s \to m_c \) should be made, and the numerical value is \( f_{\Lambda_c} = (6.4 \pm 0.7) \times 10^{-3} \text{GeV}^2 \), taken from the interval \( 1 < M_B^2 < 2 \text{GeV}^2 \) with \( s_0 \sim 10 \text{GeV}^2 \).

**B. Analysis of the LCSRs**

In the numerical analysis for the form factors, the charm quark mass is taken to be \( m_c = 1.41 \text{GeV} \) \[22\], and the other relevant parameters, \( \Lambda_c \) and \( \Lambda \) baryon masses and the value of \( |V_{cs}| \), can be found in \[23\]. We start with analysis of the twist-3 sum rules, in which only twist-3 DA is kept. Substitute the above given parameters into the LCSRs and vary the continuum threshold within the range \( s_0 = 7 - 9 \text{GeV}^2 \), we find there exist an acceptable stability in the range \( M_B^2 = 5 - 7 \text{GeV}^2 \) for the Borel parameter. The \( M_B^2 \) and the \( q^2 \) dependence for the corresponding form factors are shown in Figs. 1 and 2, respectively. Also given in Fig. 2 are two leading twist results, corresponding to only retain \( B_0 \) in the sum rules. It is apparent that in that approximation, only \( f_1 \) and \( f_2 \) survive. Apart from the leading twist DA we discuss in \[10\], there still exists another form from Chernyak, Ogloblin and Zhitnitsky \[21\]:

\[
A_{1}^{COZ}(x_i) = -21\varphi_{as}[0.52(x_1^2 + x_2^2) + 0.34x_3^2 - 2.05x_1x_2 - 0.48x_3(x_1 + x_2)], \tag{25}
\]

where \( \varphi_{as} = 120x_1x_2x_3 \) is the asymptotic DA. The corresponding result is also illustrated for contrast.

For the up to twist-6 sum rules, the stability is agreeable within the range \( s_0 = 9 - 11 \text{GeV}^2 \) and \( M_B^2 = 7 - 9 \text{GeV}^2 \). In that working region, the twist-3 contribution to \( f_1 \) is the dominant one, amounting to over 90%. However, the case is different for \( f_2 \): the main contribution comes from the twist-4 DAs and its magnitude is approximately \( \sim 1.5 \) of the twist-3 one in the whole dynamical region, but with a different sign. On account of the relatively small momentum transfer, the asymptotic behavior of DAs may not be fulfilled and we need incorporate higher conformal spin in the expansion for them. Furthermore, QCD sum rule tends to overestimate the higher conformal spin expansion.

![FIG. 1: The dependence on \( M_B^2 \) of the LCSRs for the form factors \( f_1 \) and \( f_2 \) at \( q^2 = 0 \). The continuum threshold is \( s_0 = 8 \text{GeV}^2 \) for the twist-3 result and \( s_0 = 10 \text{GeV}^2 \) for the up to twist-6 one.](image-url)
FIG. 2: The dependence on $q^2$ of the LCSRs for the form factors $f_1$ and $f_2$. The “COZ” denotes the result obtained from the COZ DAs. The continuum threshold and the Borel parameter are $s_0 = 8\text{GeV}^2$, $M_B^2 = 6\text{GeV}^2$ and $s_0 = 10\text{GeV}^2$, $M_B^2 = 8\text{GeV}^2$ for the twist-3 and the up to twist-6 results. For the two leading twist results, they are chosen to be the same with those for the twist-3 sum rule.

parameters [24], and the corresponding parameter will enter in the coefficients of the higher conformal spin expansion, which is well known as the Wandzura-Wilczek type contribution. So at the current stage, the strategy to stay with twist-3 sum rules seems to be a good choice. The $M_B^2$ and the $q^2$ dependence for the corresponding form factors are also shown in Figs. 1 and 2.

Both the form factors in the whole kinematical region, $0 < q^2 < (M_{\Lambda_c} - M)^2$, can be fitted well by the dipole formula

$$f_i(q^2) = \frac{f_i(0)}{a_2(q^2/M_{\Lambda_c}^2)^2 + a_1 q^2 / M_{\Lambda_c}^2 + 1},$$

and the uncertainties are negligible. Below in Table II we give those coefficients for two sets of parameters: $M_B^2 = 6\text{ GeV}^2$, $s_0 = 8\text{ GeV}^2$ for the twist-3 results and $M_B^2 = 8\text{ GeV}^2$, $s_0 = 10\text{ GeV}^2$ for the up to twist-6 one.

|       | twist-3  | up to twist-6 |
|-------|----------|---------------|
|       | $a_2$    | $a_1$ | $f_i(0)$       | $a_2$    | $a_1$ | $f_i(0)$       |
| $f_1$ | 1.595    | 2.203 | 0.449          | 0.993    | -1.712 | 0.392          |
| $f_2$ | 2.992    | -3.329 | -0.193         | 0.238    | -1.339 | -0.083         |

TABLE II: The dipole fit for the form factors $f_1$ and $f_2$ with $M_B^2 = 6\text{ GeV}^2$, $s_0 = 8\text{ GeV}^2$ for the twist-3 results and $M_B^2 = 8\text{ GeV}^2$, $s_0 = 10\text{ GeV}^2$ for the up to twist-6 sum rules, [16] and [17].

Using the obtained form factors, we can calculate the differential decay rate and the total decay width for the decay $\Lambda_c \to \Lambda\ell^+\nu$. The differential decay rate is shown in Fig. 3. If only the twist-3 amplitude $A_1$ is retained, we have for the total decay width $\Gamma = (7.2 \pm 2.0) \times 10^{-14}\text{GeV}$, which agrees well with the data given by the Particle Data Group [23]. This result is also in agreement with the QCD sum rule predictions made in [17]. As for the decay asymmetry parameter $\alpha$ defined in [23], we obtain $\alpha = -0.88 \pm 0.03$, which corresponds to the ratio at zero momentum transfer $f_2(0)/f_1(0) = 0.44 \pm 0.05$. That value lies very close to the recent experimental data from CLEO [26]. Note that the errors quoted above reflect the uncertainty due to the Borel parameter $M_B$ and the continuum $s_0$. The uncertainty due to the variation of the other QCD parameters is not included, which may reach 5% or more.

For a comparison, the total decay width computed from the up to twist-6 form factors is $\Gamma = (6.8 \pm 2.0) \times 10^{-14}\text{GeV}$, whose agreement with the experimental value is good, too. However, for the asymmetry parameter $\alpha$, we get $\alpha = -(0.54 \pm 0.02)$, which lies above the particle data group’s average [23] and still greater than the latest
FIG. 3: Differential decay rate for $\Lambda_c \to \Lambda \ell^+\nu$. The parameters follow those in Fig. 2.

experimental measurement [26]. The ratio of the two form factors at zero momentum transfer of that asymmetry is $f_2(0)/f_1(0) = -0.22 \pm 0.03$. This phenomenon, i.e., the twist-3 result agrees better with the experiments, may be attributed to our incomplete inclusion of the higher conformal spin components in the expansion for the DAs.

To summarize, we have given a preliminary investigation on the semileptonic decay $\Lambda_c \to \Lambda \ell^+\nu$ using the LCSR method. Sum rules for the form factors are derived and used to calculate the decay width and the asymmetry parameter. The decay width agrees well with the experimental data both for the twist-3 and the up to twist-6 sum rules, while the asymmetry’s agreement is not so good: the twist-3 DA alone can account for the experimental value but the result including higher twist contributions is not so good. This is partly due to the interplay of our incomplete inclusion of the higher conformal spin contributions and the QCD sum rule over-estimated value for $\lambda_1$.

Acknowledgments

M.Q.H. would like to thank the Abdus Salam ICTP for warm hospitality. This work was supported in part by the National Natural Science Foundation of China.

[1] I.I. Balitsky, V.M. Braun and A.V. Kolesnichenko, Nucl. Phys. B 312, 509 (1989); V.L. Chernyak and I.R. Zhitnitsky, ibid. B 345, 137 (1990).
[2] V. Braun, Light-Cone Sum Rules, [hep-ph/9801222] P. Colangelo and A. Khodjamirian, in At the Frontier of Particle Physics/Handbook of QCD, edited by M. Shifman (World Scientific, Singapore, 2001), pp. 1495-1576, [hep-ph/0010175]
[3] M. A. Shifman, A. I. Vainshtein and V. I. Zakharov, Nucl. Phys. B147, 385 (1979); B147, 448 (1979); V. A. Novikov, M. A. Shifman, A. I. Vainshtein and V. I. Zakharov, Fortschr. Phys. 32, 11 (1984).
[4] P. Ball and V. M. Braun, Phys. Rev. D 55, 5561 (1997).
[5] G. P. Lepage and S. J. Brodsky, Phys. Rev. Lett. 43, 545(1979); 43, 1625(E) (1979); G. P. Lepage and S. J. Brodsky, Phys. Rev. D. 22, 2157 (1980); V. L. Chernyak and A. R. Zhitnitsky, Phys. Rept. 112, 173 (1984).
[6] V. M. Braun, A. Lenz, N. Mahnke, and E. Stein, phys. Rev. D 65, 074011 (2002);
[7] A. Lenz, M. Wittmann and E. Stein, Phys. Lett. B 581, 199 (2004).
[8] V. M. Braun, A. Lenz and M. Wittmann, Phys. Rev. D 73, 094019 (2006).
[9] Z. G. Wang, S. L. Wan and W. M. Yang, Phys. Rev. D 73, 094011 (2006), [hep-ph/0601060].
[10] M.Q. Huang and D.W. Wang, phys. Rev. D 69, 094003 (2004).
[11] V. M. Braun, A. Lenz, G. Peters, and A.V. Radyushkin, phys. Rev. D 73, 034020 (2006).
[12] M. B. Gavela, Phys. Lett. B 83, 367 (1979).
[13] R. Perez-Marcial, R. Huerta, A. Garcia and M. Avila-Aoki, Phys. Rev. D 40, 2955 (1989); D 44, 2203(E) (1991).
[14] R. L. Singleton, Phys. Rev. D 43, 2939 (1991).
[15] H. Y. Cheng and B. Tseng, Phys. Rev. D 53, 1457 (1996); D 55, 1697(E) (1997).
[16] S. Migura, D. Merten, B. Metsch and H. R. Petry, Eur. Phys. J. A 28, 55 (2006).
[17] H. G. Dosch, E. Ferreira, M. Nielsen and R. Rosenfeld, Phys. Lett. B 431, 173 (1998); R. S. Marques de Carvalho, F. S. Navarra, M. Nielsen, E. Ferreira and H. G. Dosch, Phys. Rev. D 60, 034009 (1999).
[18] V. M. Braun, R. J. Fries, N. Mahnke, and E. Stein, Nucl. Phys. B589, 381 (2000); B607, 433(E)(2001).
[19] V. M. Braun, S. E. Derkachov, G. P. Korchemsky and A. N. Manashov, Nucl. Phys. B 553, 355 (1999).
[20] B. L. Ioffe, Nucl. Phys. B188, 317 (1981); B191, 591(E) (1981); Z. Phys. C 18, 87 (1983); V. M. Belyaev and B. Yu. Blok, Z. Phys. C 30, 151 (1983); M. A. Ivanov et al., Phys. Rev. D 61, 114010 (2000); D. W. Wang and M. Q. Huang, ibid. D 67, 074025 (2003).
[21] V. L. Chernyak, A. A. Ogloblin and I. R. Zhitnitsky, Z. Phys. C 42, 569 (1989).
[22] D. W. Wang, M. Q. Huang and C. Z. Li, Phys. Rev. D 65, 094036 (2002).
[23] S. Eidelman et al., Particle Data Group, Phys. Lett. B 592, 1 (2004).
[24] A. P. Bakulev and A. V. Radyushkin, Phys. Lett. B 271, 223 (1991); S. V. Mikhailov and A. V. Radyushkin, Phys. Rev. D 45, 1754 (1992).
[25] J. G. Korner and M. Kramer, Phys. Lett. B 275, 495 (1992); A. Kadeer, J. G. Korner and U. Moosbrugger, arXiv:hep-ph/0511019
[26] J. W. Hinson et al., CLEO Collaboration, Phys. Rev. Lett. 94, 191801 (2005).