Pion transition form factor in $k_T$ factorization

Hsiang-nan Li and Satoshi Mishima

1Institute of Physics, Academia Sinica, Taipei, Taiwan 115, Republic of China; Department of Physics, Tsing-Hua University, Hsinchu, Taiwan 300, Republic of China; Department of Physics, National Cheng-Kung University, Tainan, Taiwan 701, Republic of China; and Institute of Applied Physics, National Cheng-Chi University, Taipei, Taiwan 116, Republic of China
2Theory Group, Deutsches Elektronen-Synchrotron DESY, 22607 Hamburg, Germany

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The asymptotic and soft behaviors of the pion transition form factor $F_{\pi\gamma}(Q^2)$ involved in the process $\pi\gamma^* \rightarrow \gamma$ have been derived [1]:

$$\lim_{Q^2 \to 0} Q^2 F_{\pi\gamma}(Q^2) = \frac{\sqrt{2}}{4\pi^2 f_\pi},$$

$$\lim_{Q^2 \to \infty} Q^2 F_{\pi\gamma}(Q^2) = \frac{\sqrt{2}}{2f_\pi},$$

(1)

(2)

with $Q^2$ being the momentum transfer squared carried by the virtual photon, and $f_\pi = 0.131$ GeV is the pion decay constant. The former is predicted by perturbative QCD (PQCD) in the collinear factorization theorem [2–7], while the latter is determined from the axial anomaly in the chiral (PQCD) in the collinear factorization theorem [2–7], while the latter is determined from the axial anomaly in the chiral limit. However, the recent BABAR data on $F_{\pi\gamma}(Q^2)$ exhibits an intriguing dependence on $Q^2$ [8]: $Q^2 F_{\pi\gamma}(Q^2)$ exceeds Eq. (1) for $Q^2 > 10$ GeV$^2$, and continues to grow up to $Q^2 \approx 40$ GeV$^2$ as shown in Fig. 1. It has been commented [9] that this behavior cannot be explained by perturbative effects, like higher-order or higher-twist contributions.

It is known that one can extract nonperturbative information on the shape of the leading-twist pion distribution amplitude (DA) from the measurement of $F_{\pi\gamma}(Q^2)$ [2]. Inspired by the BABAR data, a model for the pion DA, differing from those investigated in [9], has been proposed [10],

$$\phi_\pi(x, \mu_0) = N + (1 - N)6\alpha x(1 - x),$$

(3)

with $N$ being a free constant. This model remains finite at the end points of the momentum fraction $x \to 0, 1$ for the normalization point $\mu_0 = 0.6$–0.8 GeV. It was argued [10] that the pion DAs from the instanton theory of QCD vacuum [11,12], from the Nambu-Jona-Lasinio model [13], from the chiral quark model [14], and from the large-$N_c$ Regge model [15] are expected to be rather flat. At the same time, the hard kernel, proportional to the internal quark propagator, was modified by introducing an infrared regulator $m^2$ [10]. Note that an endpoint singularity would be developed for the model in Eq. (3) without the infrared regulator. Then the factorization formula

$$Q^2 F_{\pi\gamma}(Q^2) = \frac{\sqrt{2}f_\pi}{3} \int_0^1 dx \frac{\phi_\pi(x, Q)}{x + m^2/Q^2},$$

(4)

leading to a logarithmic increase $\ln Q^2$, explains the growth of $Q^2 F_{\pi\gamma}(Q^2)$ observed by BABAR with the tuned parameters.

FIG. 1 (color online). $Q^2$ dependence of $Q^2 F_{\pi\gamma}(Q^2)$ up to LO and NLO from the PQCD approach for the asymptotic (red, dotted and dot-dashed lines, respectively) and flat (blue, dashed and solid lines, respectively) pion DAs. The points with errors are the data from CELLO [19], CLEO [20], and BABAR [8].

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terms \( N = 1.3 \pm 0.2 \) and \( m = 0.65 \pm 0.05 \) GeV. The regulator \( m^2 \) was interpreted as the inverse instanton size, which sets the scale for nonperturbative effects in a quark propagator [10].

A similar conclusion on the shape of the leading-twist pion DA has been drawn from the BABAR data, using a different formalism that takes into account the parton transverse momentum \( k_T \) [16]. The hard kernel is the same as in the collinear factorization, and the \( \ln Q^2 \) dependence comes from the integration of the pion wave function over the intrinsic parton \( k_T \) up to \( xQ \) [16]:

\[
Q^2 F_{\pi\gamma}(Q^2) = \frac{\sqrt{2} f_\pi}{6\pi} \int_0^1 dx \int_0^{xQ} k_T dk_T \psi_\pi(x, k_T),
\]

with the model

\[
\psi_\pi(x, k_T) = \frac{2\pi \phi_\pi(x)}{x(1-x)\sigma} \exp\left[-\frac{k_T^2}{2\sigma x(1-x)}\right].
\]

The cutoff of \( k_T \) at \( xQ \) guarantees that the integration over \( x \) down to 0 is finite even for a flat distribution \( \phi_\pi(x) \), with which the BABAR data are accommodated by choosing the parameter \( \sigma = 0.53 \) GeV. A concern was raised for a broad pion DA in the conventional collinear factorization [16]: the next-to-leading-order (NLO) correction to the hard kernel [17] is huge for the renormalization scale \( \Lambda_{\text{QCD}} \): the next-to-leading-order (NLO) correction to the pion DA in the conventional collinear factorization [23] is of the next-to-next-to-leading-order correction to the pion wave function, which the hard kernel [17] is huge for the renormalization scale \( \Lambda_{\text{QCD}} \).

The BABAR data are consistent with those of CELLO [19] and CLEO [20] below \( Q^2 = 9 \) GeV, from which an end point suppressed shape of the pion DA has been extracted [21,22]. Including high-power corrections, the asymptotic pion DA can also explain the CLEO data [23]. The QCD sum rule analyses of the pion DA and of the pion transition form factor led to an end point suppressed model [9,21,24–28]. It has been noticed [29] that the CLEO data favor the asymptotic form rather than the Chernyak-Zhitnitsky one [5], which emphasizes the end point region. The E791 di-jets data support the nearly asymptotic pion DA above \( M_J^2 \) = 10 GeV [30], with \( M_J^2 \) being the mass squared of the di-jets. A review on the determination of the leading-twist pion DA can be found in [31]. It seems that the measurements of the pion transition form factor at low and high \( Q^2 \) imply different shapes of the pion DA, which are contradictory to each other.

In this paper we shall study the process \( \pi^+ \pi^- \rightarrow \gamma \) using the PQCD approach [32–34] based on the \( k_T \) factorization theorem [35–37]. If a flat pion DA is favored, it will enhance the contribution from the region with small momentum fraction \( x \). Once the small \( x \) region is important, the parton transverse momentum squared \( k_T^2 \), being of the same order as \( xQ^2 \) in the internal quark propagator, is not negligible. The regulator \( m^2 \) in Eq. (4) can be interpreted as \( k_T^2 \) at LO, whose inclusion smears the end point singularity from \( x \rightarrow 0 \). Besides, the loop correction to the virtual photon vertex generates the large double logarithms \( \alpha_s \ln^2(x/k_T^2) \) and \( \alpha_s \ln x \) in the small \( x \) region [38]. The former is absorbed into the \( k_T \) dependent wave function, and organized to all orders by the \( k_T \) resummation [32,39,40]. The latter is absorbed into a jet function, and organized to all orders by the threshold resummation [41–43]. With the above resummation factors, a flat pion DA does not develop an end point singularity.

The NLO correction to the hard kernel of the pion transition form factor has been computed in the \( k_T \) factorization [38]. Fourier transforming Eq. (44) in [38] into the impact parameter \( b_T \) space, we derive

\[
F_{\pi\gamma}(Q^2) = \frac{\sqrt{2} f_\pi}{3} \int_0^1 dx \int_0^{b_T} b_T db_T \phi_\pi(x) \exp[-S(x, b_T, Q, \mu)] S_s(x, Q) K_0(\sqrt{2} Q b_T)
\]

\[
\times \left[ 1 - \frac{\alpha_s(\mu)}{4\pi} C_F \left( \ln \frac{Q^2 b_T^2}{2 s} + \gamma_E + 2 \ln x + 3 - \frac{\pi^2}{3} \right) \right],
\]

where \( C_F = 4/3 \) is a color factor and \( \gamma_E \) the Euler constant. The NLO term in the square brackets implies that the renormalization scale \( \mu \) should be chosen as \( \mu^2 \sim O(\sqrt{2} Q/b_T) \) in order to minimize the logarithm. However, we have always set \( \mu = \max(\sqrt{2} Q, 1/b_T) \) in our previous analysis of exclusive processes [33], so we shall continue to adopt this choice here. It has been confirmed that the above two choices of \( \mu \) produce almost identical results. The Sudakov exponent from the \( k_T \) resummation is given by

\[
S(x, b_T, Q, \mu) = S\left(x, \frac{Q}{\sqrt{2}}, b_T \right) + S\left((1-x), \frac{Q}{\sqrt{2}}, b_T \right)
\]

\[
+ 2 \int_{1/b_T}^\mu \frac{d\tilde{\mu}}{\tilde{\mu}} \gamma_q(\alpha_s(\tilde{\mu})),
\]

where the explicit expression of the function \( s \) and the quark anomalous dimension \( \gamma_q \) can be found in [44,45].

The power \( c = 0.3 \) in the parametrization \( S_s(x, Q) \) for the threshold resummation [46]

\[
S_s(x, Q) = 2^{1+2c} \frac{\Gamma(3/2 + c)}{\Gamma(1+c)} [x(1-x)],
\]

was derived at the scale of the B meson mass \( m_B = 5.28 \) GeV. The power-law behavior should be modified as \( Q^2 \) runs within a large range in the present case. We determine the \( Q^2 \) dependence of the parameter \( c \) by repeating the procedure in Appendix D of [46], which involves the best fit of Eq. (9) to the exact resummation formula in the Mellin space \( N \).
\[ S_c(N, Q) = \exp \left[ \frac{1}{2} \int_0^{1-N} \frac{dz}{1-z} \right. \]
\[ \times \left. \int_{(1-z)}^{(1-z)^2} \frac{d\Lambda}{\Lambda} \gamma_K(\alpha_s(\sqrt{\Lambda Q^2/2})) \right]. \]

with the anomalous dimension
\[ \gamma_K = \frac{\alpha_s}{\pi} C_F + \left( \frac{\alpha_s}{\pi} \right)^2 C_F \left[ C_A \left( \frac{67}{36} - \frac{\pi^2}{12} \right) - \frac{5}{18} n_f \right].\]

The LO and NLO predictions for the pion transition form factor from the PQCD approach are presented in Fig. 1. It is observed that the curves for \( Q^2 F_{\gamma\pi}(Q^2) \) from the asymptotic pion DA saturates quickly for \( Q^2 > 5 \text{ GeV}^2 \). That is, it is difficult to explain the continuous growth of \( Q^2 F_{\gamma\pi}(Q^2) \) above \( Q^2 = 10 \text{ GeV}^2 \) using the asymptotic model. The overshoot in the low \( Q^2 \) region is expected, since this portion of data favors an end point suppressed shape [21,22]. Note that the overshoot can be resolved within the theoretical uncertainty from varying parameters in our formalism. For example, multiplying the arguments \( xQ/\sqrt{2} \) and \((1-x)Q/\sqrt{2}\) by a factor 2 in Eq. (8) [32] leads to stronger suppression at \( x \to 0 \), and better agreement with the data at low \( Q^2 \). The curves from the flat pion DA show a good fit to the BABAR data, whose logarithmic increase with \( Q^2 \) is a combined consequence of the inclusion of the parton \( k_T \) and the employment of the flat model. We stress that the fit would deteriorate in the low \( Q^2 \) region without a careful treatment of the threshold resummation effect. To highlight this effect, we compare the results from the flat model with a constant power \( c = 0.3 \) and with a running \( c(Q^2) \) in Fig. 2. It is found that the former exceeds the data around \( Q^2 = 10 \text{ GeV}^2 \). The results from the asymptotic model are almost independent of the running of \( c(Q^2) \), because the asymptotic model decreases quickly enough at small \( x \) by itself. The NLO correction, amounting up to 20% for the flat pion DA as shown in Fig. 1, is under control in the \( k_T \) factorization. This observation is opposite to that postulated in [10], and to that obtained in the collinear factorization [16].

We should keep in mind that the BABAR data suffer from a large uncertainty in the high \( Q^2 \) region. In particular, the data point for \( Q^2 = 27.31 \text{ GeV}^2 \) drops, and is consistent with the asymptotic limit in Eq. (1). Hence, we examine the dependence of our prediction on the parameter \( \sigma \) associated with the intrinsic \( b_T \) dependence in Eq. (13). The results displayed in Fig. 3(a) indicate that the slope of the

### Table 1

| \( Q \) (GeV) | 2   | 3   | 4   | 5   | 6   |
|--------------|-----|-----|-----|-----|-----|
| \( c \)      | 1.7 | 0.75| 0.47| 0.32| 0.25|

FIG. 2 (color online). \( Q^2 \) dependence of \( Q^2 F_{\gamma\pi}(Q^2) \) up to NLO for the flat model with the power \( c = 0.3 \) (red, dashed line) and \( c(Q^2) \) (blue, solid line) in Eq. (12).
with the Gegenbauer polynomial mentioned [10] that a very broad pion DA, vanishing at the end points but with a rapid increase \[a_2 \approx 0.8\] especially in the high \(Q^2\) region, and saturates around \(Q^2 = 10 \text{ GeV}^2\). Only with an extremely large \(a_2 \approx 0.8\), which also emphasizes the end point regions, the prediction becomes close to that from the flat pion DA. It has been mentioned [10] that a very broad pion DA, vanishing at the end points but with a rapid increase \([\phi'_q(0) / 6 \gg 1]\), serves the purpose of accommodating the BABAR data. However, the big value \(a_2 \approx 0.8\) may render the Gegenbauer expansion starting with the asymptotic model questionable, as claimed in [10]. In summary, the precise measurement of the pion transition form factor in the high \(Q^2\) region will settle the issue on the shape of the leading-twist pion DA: if the growth with \(Q^2\) observed recently persists, the flat model is favored. If the growth becomes milder in the future, the asymptotic model with nonvanishing higher Gegenbauer terms will still be allowed.

Finally, we comment on the \(k_T\) factorization formalism developed in [42], which involves on-shell partons in the calculation of hard kernels. The usual procedure of factorizing a \(k_T\) dependent wave function is to neglect the minus component \(k\) in a hard kernel, and then to integrate out \(k\) in the general wave function that depends on the four components \(k^+ = xp^+, k^-,\) and \(k_T\),

\[
\psi(x, k_T) = \int dk^- \Psi(x, k^-, k_T).
\] (16)

This is the reason why a parton, participating in the hard scattering, carries an off-shell momentum \((k^+, 0, k_T)\) in our formalism [49]. The prescription for calculating a \(k_T\) dependent hard kernel, the definition for a \(k_T\) dependent wave function, and the gauge invariance of the \(k_T\) factorization have been carefully addressed in [38,49], which confront the criticism from [50,51]. The transverse momentum squared \(k_T^2\) appears in the hard kernel for the pion transition form factor through the internal quark invariant mass, \((k - P_\gamma)^2 = -xQ^2 - k_T^2\), for the final-state photon momentum \(P_\gamma = (0, P^-, 0)\) and \(k^2 = k_T^2\). Assuming on-shell partons, the hard kernel is independent of \(k_T^2\) because of \((k - P_\gamma)^2 = -xQ^2\) for \(k^2 = 0\). Therefore, it will be difficult to explain the BABAR data using the formalism of [42].

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1The authors of [51] rebutted our comments addressed in [49] on [50]. They claimed that the \(k_T\) factorization theorem for exclusive processes violates gauge invariance for the following reasons: (1) Our method for the one-loop calculation of Fig. 2(b) in [49] leads to a wave function nonvanishing for \(x > 1\), and violates the translational invariance, since the contribution from \(q^+ < 0\) is nonzero, where \(q\) is the loop momentum. (2) Our method depends on the contour chosen in the \(q^-\) integration. We disagree with both statements. In our method, the \(q^- < 0\) contribution comes only from the limit \(q^+ \rightarrow 0\) from the left, which corresponds to \(x \rightarrow 1\), and does not violate the translational invariance. As for the latter, the authors of [51] missed the point of including the semicircles, whose purpose is to avoid the ambiguity from \(q^+ \rightarrow 0\), when the double pole \(q^+ = (q_1^- - i\epsilon)/(2q^+)\) moves to infinity. If choosing the lower semicircle, the poles will cross the semicircle as \(q^+ \rightarrow 0\), and the ambiguity comes back. Besides, the authors of [51] relied on the assumption that all contributions from the semicircles in the \(q^+\) contour integration vanish in the limit of the semicircle radius \(R \rightarrow \infty\), but it is incorrect. As explained in [49], these contributions do not vanish as \(q^+ \rightarrow 0\). Note that the same problem has been discussed in literature, e.g. in [52]. Moreover, the reply in [51] is selective; e.g. there was no answer to our question raised in [49]: why is the singularity from a highly off-shell gluon with \(q^2 \rightarrow \infty\) a light-cone singularity?
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