Transverse momentum dependence in the 
perturbative 
calculation of pion form factor

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Abstract

By reanalysing transverse momentum dependence in the perturbative calculation of pion form factor an improved expression of pion form factor which takes into account the transverse momentum dependence in hard scattering amplitude and intrinsic transverse momentum dependence associated with pion wave functions is given to leading order, which is available for momentum transfers of the order of a few GeV as well as for $Q \to \infty$. Our scheme can be extended to evaluate the contributions to the pion form factor beyond leading order.

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1 Introduction

Exclusive processes in perturbative quantum chromodynamics (PQCD) were first studied by Brodsky and Lepage [1] many years ago. Recently, the studies of the exclusive processes at experimentally accessible momentum transfers in the framework of PQCD have received much interest. The statement [2, 3] that the applicability of PQCD to the exclusive processes at experimentally accessible momentum transfers is questionable has been challenged [4-11]. Huang and Shen [4] pointed out that the applicability of PQCD to form factor is questionable only as momentum transfer $Q^2 \leq 4 \text{ GeV}^2$ through the study on the pion form factor by reanalyzing the contributions from end-point regions. Botts and Sterman [5] proposed a formulation in which both Sudakov and nonleading logarithmic corrections to independent (Landshoff) scatterings [12] of valence quarks are organized systematically. Li and Sterman [6] gave out a modified expression for the pion form factor by taking into account the customarily neglected partonic transverse momentum as well as Sudakov correction. They reached a similar conclusion as [4]: PQCD begins to be self-consistent at about $Q \sim 20\Lambda_{QCD}$. Jakob and Kroll [8] demonstrated that for momentum transfers of the order of a few GeV the intrinsic transverse momentum dependence of wave function leads to a substantial suppression of the perturbative contributions, which should be considered besides Sudakov suppression. The base of the most previous discussions [5-11] is the formalism in Ref. [5] which is suitable for studying the large $Q$ region since it sets $b \to 0$ in the integral of the wave function in respecting the intrinsic transverse momentum dependence of wave function. In this paper we re-analyse the PQCD calculation for the pion form factor at momentum transfers of the order of a few GeV. We will give out an improved expression for the pion form factor which takes into account the transverse momentum dependence in the gluon propagator as well as in the fermion propagator in the hard scattering amplitude $T_H$ and the intrinsic transverse momentum dependence associated with the pion wave function.
This expression is available for momentum transfers of the order of a few GeV as well as for $Q \rightarrow \infty$ since it does not make the approximation $b \rightarrow 0$. The formalism of Ref. [8] is just an approximate expression in respecting the intrinsic transverse momentum dependence. This approximation brings sizeable effect on the numerical prediction for the pion form factor in the momentum transfer $Q \sim$ a few GeV region. It is also found that the transverse momentum dependence of the fermion propagator in $T_H$ leads to a mild reduction of the prediction for the pion form factor in the same momentum transfer region. The remainder of the paper is organized as follows. Sect. 2 reviews Li-Sterman’s formalism for the pion form factor. In Sect. 3, we discuss the transverse momentum dependence in pion wave function beside the one in the hard scattering amplitude $T_H$ in perturbative calculation. In Sect. 4, we do our numerical calculations. Finally, in Sect. 5, we give a summary.

2 Brief review of Li-Sterman’s formalism

Taking into account the transverse momenta $k_T$ that flow from the wave functions through the hard scattering leads to a factorization form with two wave functions $\psi(x_i, k_{T_i})$ corresponding to the external pions, combined with a new hard-scattering function $T_H(x_1, x_2, Q, k_{T_1}, k_{T_2})$, which depends in general on transverse as well as longitudinal momenta [6],

$$F_\pi(Q^2) = \int_0^1 dx_1 dx_2 \int \frac{d^2k_{T_1}}{16\pi^3} \frac{d^2k_{T_2}}{16\pi^3} \psi(x_1, k_{T_1}) \times T_H(x_1, x_2, Q, k_{T_1}, k_{T_2}) \psi(x_2, k_{T_2}).$$ (1)

In this form, both soft and collinear logarithmic enhancements are factorized into the functions $\psi$. At the lowest order, $T_H$ reads \footnote{A frame with $q_T = 0$ has been adopted to obtain this expression of $T_H$ in [8]. The expressions of $T_H$ with different momentum assignment will be analysed in detail [13].}

$$T_H(x_1, x_2, Q, k_{T_1}, k_{T_2}) = \frac{16\pi C_F \alpha_s(\mu) x_1 Q^2}{(x_1 Q^2 + k_{T_1}^2)(x_1 x_2 Q^2 + (k_{T_1} + k_{T_2})^2)}.\quad (2)$$
Neglecting the transverse momentum dependence in the fermion propagator, Eq. (2) becomes

\[ T_H(x_1, x_2, Q, k_{T_1}, k_{T_2}) = \frac{16\pi C_F\alpha_s(\mu)}{x_1 x_2 Q^2 + (k_{T_1} + k_{T_2})^2}. \]  

(3)

The next step is to re-express Eq. (1) in terms of the Fourier transformation variables in the transverse configuration space [6]. Observing that \( T_H \) in Eq. (3) depends on only a single combination of the transverse momenta \( (k_{T_1} + k_{T_2}) \), the Fourier transformation of Eq. (1) involves only a single integral of Fourier transform variable

\[ F_\pi(Q^2) = \int_0^1 dx_1 dx_2 \frac{db}{(2\pi)^2} \varphi(x_1, b, \mu) T_H(x_1, x_2, Q, b, \mu) \varphi(x_2, b, \mu). \]  

(4)

The wave functions in \( b \)-space, \( \varphi(x, b, \mu) \) take into account an infinite summation of higher-order effects associated with the elastic scattering of valence partons, which gives out Sudakov suppression to the large-\( b \) and small-\( x \) regions.

The asymptotic behavior of \( \varphi(x, b, \mu) \) at large \( Q^2 \) has been obtained in Ref. [5]

\[ \varphi(x, b, \mu) = \exp \left[ -s(x, b, Q) - s(1 - x, b, Q) - 2 \int_{1/b}^\mu \frac{d\hat{\mu}}{\hat{\mu}} \gamma_q(g(\hat{\mu})) \right] \times \phi \left( x, \frac{1}{b} \right), \]  

(5)

where \( \gamma_q = -\alpha_s/\pi \) is the quark anomalous dimension in the axial gauge. \( s(\xi, b, Q) \) is Sudakov exponential factor, which reads [3, 11]

\[
s(\xi, b, Q) = \frac{A^{(1)}(1)}{2\beta_1} \hat{q} \ln \left( \frac{\hat{q}}{-b} \right) + \frac{A^{(2)}}{4\beta_1^2} \left( \frac{\hat{q}}{-b} - 1 \right) - \frac{A^{(1)}}{2\beta_1} (\hat{q} + \hat{b}) \\
- \frac{A^{(1)}\beta_2}{4\beta_1^3} \hat{q} \left[ \ln(-2\hat{b}) + 1 - \ln(-2\hat{q}) + 1 \right] \\
- \left( \frac{A^{(2)}}{4\beta_1^2} - \frac{A^{(1)}(1)}{4\beta_1} \ln\left( \frac{1}{2e^{2\gamma-1}} \right) \right) \ln \left( \frac{\hat{q}}{-b} \right) \\
+ \frac{A^{(1)}\beta_2}{8\beta_1^3} \left[ \ln^2(2\hat{q}) - \ln^2(-2\hat{b}) \right], \tag{6}
\]

where

\[
\hat{q} = \ln[\xi Q/(\sqrt{2}\Lambda)], \quad \hat{b} = \ln(b\Lambda), \\
\beta_1 = \frac{33 - 2n_f}{12}, \quad \beta_2 = \frac{153 - 19n_f}{24}, \\
A^{(1)} = \frac{4}{3}, \quad A^{(2)} = \frac{67}{9} - \frac{1}{3}\pi^2 - \frac{10}{27}n_f + \frac{8}{3}\beta_1 \ln\left( \frac{1}{2e^{\gamma}} \right). \tag{7}
\]
\(n_f\) is the number of quark flavors. \(\gamma\) is the Euler constant. The factor \(e^{-s(\xi,b,Q)}\) which induces an enhancement in small-\(b\) regions has been set to unity whenever \(\xi \leq \sqrt{2}/(bQ)\). The coefficients in Eq. (8) are different from Ref. [3] since there are some algebraically mistakes in the previous expression of \(s(\xi,b,Q)\), which has been pointed in Ref. [11]. \(\phi(x,1/b)\) in Eq. (5) is a “soft” wave function calculated with gluons of transverse momentum \(k_T \leq 1/b\),

\[
\phi(x,1/b) = \int_{k_T \leq 1/b} \frac{d^2 k_T}{16\pi^3} \psi(x,k_T).
\]  

(8)

Neglecting the \(b\)-dependence of the function \(\phi(x,1/b)\) and performing Fourier transformation for Eq. (3), the expression Eq. (1) becomes [3]

\[
F_\pi(Q^2) = 16\pi C_F \int_0^1 dx_1 dx_2 \int_0^\infty db \alpha_s(t) K_0(\sqrt{x_1 x_2 Qb}) \times \exp (-S(x_1,x_2,Q,b,t)) \phi(x_1) \phi(x_2),
\]  

(9)

where \(\phi(x)\) is defined as

\[
\phi(x) = \int \frac{d^2 k_T}{16\pi^3} \psi(x,k_T),
\]  

(10)

\(K_0\) is the modified Bessel function of order zero, and

\[
S(x_1,x_2,Q,b,t) = \left[ \sum_{i=1}^2 (s(x_i,b,Q) + s(1-x_i,b,Q)) - \frac{2}{\beta_1} \ln \left( \frac{\hat{t}}{-\hat{b}} \right) \right].
\]  

(11)

3 \(k_T\)-dependence in perturbative calculation

Following Ref. [3], we find that \(\phi(x,1/b)\) in Eq. (5) should be replaced by the amplitude

\[
\varphi^{(1/b)}(x,b) = \int_{k_T \leq 1/b} \frac{d^2 k_T}{16\pi^3} e^{ik_T \cdot b} \psi(x,k_T).
\]  

(12)

Through the requirement that \(e^{ik_T \cdot b} \approx 1\), as \(Q \to \infty\) \((b \to 0)\). Ref. [3] proposed \(\varphi^{(1/b)}(x,b)\) can be expressed approximately by \(\phi(x,1/b)\). For simplicity, Ref. [3] neglected the \(b\)-dependence of the function \(\varphi^{(1/b)}(x,b)\), namely neglected the intrinsic
transverse momentum dependence of wave function. Then \( \varphi^{(1/b)}(x, b) \) can be pulled out of the \( b \)-integral. Jakob and Kroll \[8\] improved this approximation and proposed to take into account the intrinsic transverse momentum dependence by a function \( \psi(x, b) \) which is the Fourier transformation of wave function,

\[
\psi(x, b) = \int_{k_T \leq \infty} \frac{d^2 k_T}{16\pi^3} e^{ik_T \cdot b} \varphi(x, k_T),
\]

(13)

It can be seen that \( \psi(x, b) \) is just an approximate expression for \( \varphi^{(1/b)}(x, b) \), i.e. \( \psi(x, b) = \varphi^{(\infty)}(x, b) \). When \( b \to 0 \) (namely \( Q \to \infty \)) they are consistent with each other. When \( Q \) is the order of a few GeV, the difference may be sizable. The difference between \( \varphi^{(1/b)}(x, b) \) and \( \psi(x, b) \) can be expressed by a function \( D(x, b) \),

\[
D(x, b) = \int_{1/b}^{\infty} \frac{d^2 k_T}{16\pi^3} e^{ik_T \cdot b} \psi(x, k_T),
\]

(14)

which increases as \( b \) becomes large. Eq. (13) enlarges the upper limit of the integral Eq. (12) from \( 1/b \) to \( \infty \), which corresponds to evaluate the contributions from the perturbative tail of wave function once again. Sudakov form factor provides much more suppression for large \( Q \) than for small \( Q \). Thus substituting \( \psi(x, b) \) for \( \varphi^{(1/b)}(x, b) \) does not effect the pion form factor for large \( Q \). As \( Q \) is the order of a few GeV one should investigate the effects of this substitution.

The Fourier-transformed hard-scattering amplitude from Eq. (3) reads

\[
T_H(x_1, x_2, Q, b, \mu) = 16\pi\alpha_s(\mu) C_F K_0(\sqrt{x_1 x_2} Q b).
\]

(15)

The renormalization group applied to \( T_H \) gives

\[
T_H(x_1, x_2, Q, b, \mu) = \exp \left[ -4 \int_{\mu}^{t} \frac{d\bar{\mu}}{\bar{\mu}} \gamma_q(g(\bar{\mu})) \right] \times T_H(x_1, x_2, Q, b, t).
\]

(16)

The variable \( t \) was taken as \( t = \max(\sqrt{x_1 x_2} Q, 1/b) \) and a cut-off on the running coupling constant (\( \alpha_s \leq 0.7 \)) for large-\( b \) region was made to guarantee PQCD to be self-consistent in Ref. [6]. However, in the regions of small \( x_1 x_2 Q^2 \) and large \( b \), the nonperturbative contributions maybe important. For example, the multi-gluon
exchange can occur between quark and antiquark and the transverse momentum intrinsic to the bound state wave-functions flows through all the propagators. Ref. [14] suggested a frozen $\alpha_s$ to take into account these effects. Instead of the cut-off of $\alpha_s$ in the $b$-space, Ref. [11] suggested that the coupling constant is frozen at $b \sim 1/(\sqrt{\langle k_T^2 \rangle})$ and the variable $t$ is taken as

$$t = \max \left( \sqrt{x_1x_2Q}, 1/b_F \right),$$

where

$$b_F = \begin{cases} 
  b & \text{if } 1/b \geq \sqrt{\langle k_T^2 \rangle} \\
  1/\sqrt{\langle k_T^2 \rangle} & \text{if } 1/b < \sqrt{\langle k_T^2 \rangle},
\end{cases}$$

and $\sqrt{\langle k_T^2 \rangle}$ is the average transverse momentum of the pion. In this way, the perturbative contributions to the pion form factor can be calculated from the present energy with a reasonable $\alpha_s$. It should be emphasized that the average transverse momentum $\sqrt{\langle k_T^2 \rangle}$ is determined definitely by the hadronic wave function.

Combining Eqs. (4), (5), (12), (15) and (16), we have

$$F_{\pi}(Q^2) = 16\pi C_F \int_0^1 dx_1 dx_2 \int_0^\infty b \, db \alpha_s(t) K_0(\sqrt{x_1x_2Qb}) \times \exp \left( -S(x_1, x_2, Q, b, t) \right) \varphi^{(1/b)}(x_1, b)\varphi^{(1/b)}(x_2, b),$$

where

$$S(x_1, x_2, Q, b, t) = \left[ \sum_{i=1}^2 \left( s(x_i, b, Q) + s(1-x_i, b, Q) \right) - \frac{2}{\beta_1} \ln \frac{\hat{t}}{-\hat{b}} \right].$$

Eq. (19) is an improved expression for the pion form factor. From our formalism, it can be found that: (a) neglecting the transverse momentum dependence associated with the wave function, Eq. (19) becomes the expression in Ref. [6] (See Eq. (9)); (b) approximating $\varphi^{(1/b)}(x, b)$ with $\psi(x, b)$, Eq. (19) becomes the expression in Ref. [8],

$$F_{\pi}(Q^2) = 16\pi C_F \int_0^1 dx_1 dx_2 \int_0^\infty b \, db \alpha_s(t) K_0(\sqrt{x_1x_2Qb}) \times \exp \left( -S(x_1, x_2, Q, b, t) \right) \psi(x_1, b)\psi(x_2, b).$$
where $\psi(x, b)$ is the Fourier transformation of the wave function (Eq. (13)).

The $k_T$-dependence in the fermion propagator contributes to $T_H$ a factor $\frac{x_1 Q^2}{x_1 Q^2 + k_{T_1}^2}$ which involves only a single transverse momentum corresponding to one in the external pion (see Eqs. (2) and (3)) and this factor makes the Fourier transformation for $T_H$ involve multi-$b$-integrals. Although the $k_T$-dependence in this factor is linear rather than quadratic in the $x$'s, it may bring some effects at the end-point region $x_1 \sim 0$. We can consider this factor and keep the simplification of the Fourier transformation for $T_H$ at the same time by combining it with the corresponding wave function to define a new “wave function”,

$$\tilde{\psi}(x_1, k_{T_1}) = \frac{x_1 Q^2}{x_1 Q^2 + k_{T_1}^2} \psi(x_1, k_{T_1}).$$

Substituting $\tilde{\psi}$ for $\psi$ in Eq. (1), the transverse momentum dependence of the fermion propagator in $T_H$ can be taken into account easily in the perturbative calculations for the pion form factor.

4 Numerical calculations

We adopt two models of the pion wave function: (a) the BHL wave function \[14\]

$$\psi^{(a)}(x, k_T) = A \exp \left[ -\frac{k_T^2 + m^2}{8\beta^2 x(1-x)} \right],$$

where \[15\] $A = 32 \text{ GeV}^{-1}$, $\beta = 0.385 \text{ GeV}$, $m = 289 \text{ MeV}$ and $\sqrt{\langle k_T^2 \rangle} = 356 \text{ MeV}$; (b) the CZ-like wave function \[16-18\]

$$\psi^{(b)}(x, k_T) = A (1 - 2x)^2 \exp \left[ -\frac{k_T^2 + m^2}{8\beta^2 x(1-x)} \right],$$

where \[15\] $A = 136 \text{ GeV}^{-1}$, $\beta = 0.455 \text{ GeV}$, $m = 342 \text{ MeV}$ and $\sqrt{\langle k_T^2 \rangle} = 343 \text{ MeV}$.

In order to discuss the effect of transverse momentum dependence associated with wave function, first, we neglect the $k_T$-dependence of the fermion propagator in $T_H$. Fig. 1 compares the behaviors of functions $\varphi^{(1/b)}(x, b)$, $\psi(x, b)$ and $\phi(x, 1/b)$.
All of them suppress the contributions from large-$b$ region, but the suppression behaviors are different in quantity. In the $b \sim 0$ region, $\phi(x, 1/b)$ and $\psi(x, b)$ are good approximation for $\varphi^{(1/b)}(x, b)$, while in the large-$b$ region that is questionable. Sudakov form factor suppresses the contributions from large-$b$ region in a mild way for small momentum transfer $Q$ than for large $Q$. Thus it is not a good approximation with $\psi(x, 1/b)$ to respect the intrinsic $k_T$-dependence of the wave functions for small $Q$, since much more contributions come from large-$b$ region. Numerical evaluations of $F_\pi$ through Eqs. (9), (19) and (21), using the BHL and CZ-like wave functions, which are plotted in Figs. 2 and 3 confirm the observation made in Fig. 1. Comparing with the original PQCD prediction[1] of the pion form factor (namely neglecting the $k_T$-dependence in $T_H$ as well as in wave functions; the dotted line), Sudakov form factor gives a suppression effect to $F_\pi$, which was first pointed out by Brodsky and Lepage [1]. Comparing with the result obtained from Eq.(9) (namely neglecting the intrinsic $k_T$-dependence of wave function; the dash-dotted line), $\varphi^{1/b}(x, b)$ (the solid line) suppresses the contribution from PQCD by about 50%, while $\psi(x, b)$ (the dashed line) suppresses by about 30% at $Q = 2\text{GeV}$ for the BHL wave function. The corresponding quantities are 55% and 35% respectively in the case of the CZ-like wave function. The intrinsic $k_T$-dependence of the wave function provides additional substantial suppression for $F_\pi$ besides Sudakov form factor. $\varphi^{1/b}(x, b)$ suppresses the contribution from PQCD more strongly than $\psi(x, b)$ does. The numerical predictions of $F_\pi$ obtained from $\varphi^{1/b}(x, b)$ and $\psi(x, b)$ are different in the momentum transfer $Q \sim a$ few GeV region.

We evaluate the effect of the $k_T$-dependence in the fermion propagator also. It leads to a small reduction of the prediction for $F_\pi$ by about 10% for $Q \sim 3\text{GeV}$, which is coincides with Ref. [1].
5 Summary

In this paper we re-analyze the transverse momentum dependence in the perturbative calculation of pion form factor at momentum transfers of the order of a few GeV. We give out an improved expression for the pion form factor which takes into account gluon propagator as well as fermion propagator transverse momentum dependence in the hard scattering amplitude and intrinsic transverse momentum dependence associated with pion wave functions. It is found that the previous approach is just an approximate expression in respecting the transverse momentum dependence associated with wave functions. This approximation brings sizeable effect on the numerical predictions for the pion form factor in the momentum transfer $Q \sim$ a few GeV region. It is also found that the transverse momentum dependence of fermion propagator in $T_H$ leads to a mild reduction of the prediction for the pion form factor in the same momentum transfer region.

We would like point out one more times that our formalism is available for momentum transfer of the order of a few GeV as well as for $Q \to \infty$, and our scheme can be extended to evaluate the higher order and higher helicity contributions for the pion form factor. The more studies on the $k_T$-dependence in the hard scattering amplitude $T_H$ are in proceeding.

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Figure Captions

**Fig. 1.** The functions $\varphi^{(1/b)}(x, b)$ (solid line), $\psi(x, 1/b)$ (dashed line) and $\phi(x, b)$ (dotted line) which are adopted to respect the intrinsic transverse momentum dependence.

$x = 0.5$, $Q^2 = 4 \text{ GeV}^2$, for the BHL wave function.

**Fig. 2.** The pion form factor calculated with the BHL wave function. The solid and dashed lines are for $\varphi^{(1/b)}(x, b)$ and $\psi(x, 1/b)$ respectively. The dash-dotted line is obtained by neglecting the intrinsic transverse momentum dependence in wave functions. The dotted line is evaluated by neglecting $k_T$-dependence in both $T_H$ and wave functions.

**Fig. 3.** Similar to Fig. 2. The pion form factor calculated with the CZ-like wave function.
Fig. 2

$F_2(Q^2) \ (GeV^2)$

$Q^2 \ (GeV^2)$
Fig. 3
Fig. 1