Abstract

We calculate electromagnetic pion form factors with an analytic model for $\alpha_s(Q^2)$ which is infrared (IR) finite without invoking a “freezing” hypothesis. We show that for the asymptotic pion distribution amplitude, $F_{\pi\gamma\gamma}$ agrees well with the data, whereas the IR-enhanced hard contribution to $F_\pi$ and the soft (nonfactorizing) part can jointly account for the data.

12.38.Bx, 12.38.Cy, 12.38.Lg, 13.40.Gp
The issue of computing exclusive processes, like electromagnetic form factors, within QCD is of fundamental interest because it reveals the basic structure of hadrons. But in contrast to inclusive processes, there is much uncertainty about the applicability of perturbative QCD at laboratory momenta. It seems that the more interpretations this subject inspires and grounds the more unsettled it is.

This paper attempts to bring together some crucial results about pion form factors and combine them with novel theoretical developments concerning the infrared (IR) regime of QCD, in order to benchmark the status of our current understanding. Such an approach appears attractive since the incorporation of nonperturbative power corrections in the perturbative domain may improve both the IR insensitivity of exclusive observables and the self-consistency of calculations entrusted to perturbative QCD.

To this end, the conventional representation of the running strong coupling “constant” is given up in favor of an analytic model, recently proposed by Shirkov and Solovtsov, which incorporates a single power correction to remove the Landau singularity. Bearing in mind that the definition of $\alpha_s$ beyond two loops cannot be uniquely fixed, one may regard the ambiguity in the IR modification of the running coupling as resembling the freedom of adopting a particular non-IR-finite renormalization scheme.

The Shirkov-Solovtsov model employs Lehmann analyticity to bridge the regions of small and large momenta by changing the one-loop effective coupling to read

$$\tilde{\alpha}_s^{(1)}(Q^2) = \frac{4\pi}{\beta_0} \left[ \frac{1}{\ln(Q^2/\Lambda^2)} + \frac{\Lambda^2}{\Lambda^2 - Q^2} \right],$$

where $\Lambda \equiv \Lambda_{\text{QCD}}$ is the QCD scale parameter. This approach provides a nonperturbative regularization at low momenta and leads to a universal value of the coupling constant at zero momentum $\tilde{\alpha}_s^{(1)}(Q^2 = 0) = 4\pi/\beta_0 \approx 1.396$ (for three flavors), defined only by group constants, i.e., avoiding the introduction of external parameters, like an (effective) gluon mass to “freeze” the coupling at low momentum scales.

Note that this limiting value (i) does not depend on the scale parameter $\Lambda$ – this being a consequence of renormalization group invariance – and (ii) extends to the two-loop order, and beyond, i.e., $\tilde{\alpha}_s^{(2)}(Q^2 = 0) = \tilde{\alpha}_s^{(1)}(Q^2 = 0) \equiv \tilde{\alpha}_s(Q^2 = 0)$. (In the following the bar is dropped.) Hence, in contrast to standard perturbation theory, the IR limit of the coupling constant is stable, i.e., does not depend on higher-order corrections, and is therefore universal. As a result, the running coupling also shows IR stability. This is tightly connected to the nonperturbative contribution $\propto \exp(-4\pi/\beta_0)$ which ensures analytic behavior in the IR domain by eliminating the ghost pole at $Q^2 = \Lambda_{\text{QCD}}^2$.

At very low momentum values, say, below 1 GeV, $\Lambda_{\text{QCD}}$ in this model deviates from that used in minimal subtraction schemes. However, since we are primarily interested in a region of momenta which is much larger than this scale, the role of this renormalization-scheme dependence is only marginal. In our investigation we use $\Lambda_{\text{QCD}}^{\text{an}(3)} = 242$ MeV which corresponds to $\Lambda_{\text{QCD}}^{\text{MS}(3)} = 200$ MeV.

This analytic model for the strong running coupling is very suitable for calculations of exclusive amplitudes, mainly for two reasons: Firstly, it ensures IR safety of the factorized short-distance part without invoking the additional assumption of saturation of color forces by using a gluon mass – extensively used up to now in form-factor calculations (see, for example, [12]). Furthermore, the Sudakov form factor [13] does not have to serve as an...
IR protector against $\alpha_s$ singularities. Hence, the extra constraint of using the maximum between the longitudinal and the transverse scale, as argument of $\alpha_s$, proposed in [14], becomes superfluous. This is a serious advantage relative to previous analyses because now one is able to choose the unphysical constants [15], which parametrize different factorization and renormalization schemes, in such a way as to optimize calculated observables (for a more detailed discussion of this point, we refer to [11]). Second, and more important, the non-logarithmic term in Eq. (1) enters all anomalous dimensions, viz. the cusp anomalous dimension, which gives rise to the Sudakov form factor [16,17,15,14,20], as well as the quark anomalous dimension which governs evolution. As a result, the suppression due to transverse momenta, intrinsic [21] and those generated by radiative corrections [14], is counteracted, and hence there is no reduction of the form-factor magnitude.

We are going to show in this work that the enhancement effect originating from the power correction to the running coupling is enough for the asymptotic pion distribution amplitude to contribute (at leading order) to the spacelike electromagnetic form factor of the pion, $F_{\pi}(Q^2)$, a hard part that can account for almost half of the form-factor magnitude relative to the existing data [22,23].

On the other hand, the transition form factor, $F_{\pi\gamma^*\gamma}(Q^2)$, is only slightly changed, as compared to the result given in [24], and matches the recent high-precision CLEO data [25] as good as the dipole fit. Also the older CELLO data [26] at lower $Q^2$ are well reproduced (see below).

In both cases, no adjustment of the theoretical predictions to the experimental data is involved.

Therefore, there appears to be no need to reanimate endpoint-concentrated pion distribution amplitudes, proposed by Chernyak and Zhitnitsky (CZ) [27], in order to make contact with the experimental data – as recently attempted in [28]. We are less enthusiastic about using such distribution amplitudes because of the following serious disadvantages:

(i) It has been recently shown [24,29] that distribution amplitudes of the CZ-type lead to a $\pi\gamma$ transition form factor which significantly overestimates the CLEO data just mentioned. Our reasons for skepticism parallel the arguments given in [24] and will not be repeated here. The excellent agreement between theory (QCD) and measurement for this process, already at leading order, when the asymptotic pion wave function is used, cannot be overemphasized. Note in this context that the calculation of Cao et al. [30], which predicts for the CZ wave function smaller values of $F_{\pi\gamma^*\gamma}(Q^2)$ than the data, uses for modeling the $k_\perp$ distribution in the pion an ansatz that strongly suppresses the endpoint region. Hence, in effect, their wave function, though claimed to be the CZ one, excludes this region and yields therefore a result even smaller than the one predicted by the asymptotic distribution amplitude in the range of $Q^2$ where there are data. Independently, investigations [31,32] based on QCD sum rules come to a comparably good description of $F_{\pi\gamma^*\gamma}(Q^2)$ at not too low $Q^2$ on the basis of local duality without presuming the asymptotic form of the pion distribution amplitude, but favoring again a shape close to that.

(ii) Distribution amplitudes of the CZ-type yield a direct-overlap, i.e., soft contribution, to $F_{\pi}(Q^2)$ which turns out to be of the same large order of magnitude [33,34,35,24] as that resulting from the convolution with the hard-scattering amplitude [30,17,27], or is even larger, at currently probed $Q^2$ values. Inclusion of this contribution into the pion form factor leads eventually to a total result which overestimates the existing data considerably – even
allowing for some double counting of hard and soft transverse momenta near the transition region. Though the present quality of the high-$Q^2$ data $^{22,23}$ on the spacelike pion form factor is quite poor, the trend seems to be indicative.

(iii) The underlying QCD sum rules analysis of $^{27}$ suffers with respect to stability – as outlined by Radyushkin $^{34}$. As a result, the duality interval increases with moment order $N$, meaning that the (nonperturbative) condensate contributions grow with order relative to the perturbative term. A direct consequence of this is that the moments of the pion distribution amplitude for $N = 2, 4$, extracted from these sum rules, are artificially enhanced. Such large moment values can only be reproduced by a double-humped endpoint-concentrated distribution amplitude and correspond to the basic assumption that vacuum field fluctuations have infinite size, or equivalently that vacuum quarks have exactly zero virtuality $^{34}$.

(iv) The characteristic humps in the endpoint regions ($x = 0, 1$) are not generic, but merely the result of truncating the eigenfunctions expansion of Gegenbauer polynomials at polynomial order two while keeping the normalization fixed to unity. Including higher and higher order polynomials, the humps become less and less prominent and the central region ($x = 1/2$) gets enhanced. To this point, we mention that an independent QCD sum rules analysis $^{38}$ gives the constraint $\phi_{\pi}(x = 1/2) = 1.2 \pm 0.3$, which is close to the value $\phi_{as}(x = 1/2) = 3/2$, and definitely violated by the CZ amplitude.

Physically, the source of the endpoint enhancement of CZ-type distribution amplitudes can be understood as follows. If the vacuum quark virtuality is zero, an infinite number of such quanta can migrate from the vacuum to the pion state at zero energy cost. This happens exactly in the kinematic region $x = 0$ or $\bar{x} = 1 - x = 0$ and leads to a strong enhancement of that region at the expense of depleting the amplitude for finding configurations in which the quark and the antiquark, or more precisely, the struck and the spectator partons, share almost equal momentum fractions around $x = 1/2$. In the pion, configurations close to the kinematic boundary contain one leading parton, which picks up almost all of the injected momentum, and an infinite number of wee partons $\sim 1/x$ with no definite transverse positions relative to the electromagnetic probe, which constitute a soft cloud. In this regime, gluons have very small virtualities and therefore it is inconsistent to assume hard-gluon rescattering, i.e., the factorization of a short-distance part in the exclusive amplitude becomes invalid. This region of momenta has to be treated separately on the basis of the Feynman mechanism just described, but a theoretical approach from first principles, though of paramount importance, is still lacking. On the other hand, if the vacuum virtuality is sizeable, say, of the order of $\Lambda_{QCD}$ or even larger $^{39}$, then an energy gap might exist that prohibits the diffusion of vacuum quarks into the pion state, and hence Feynman-type configurations are insulated from those for which hard-gluon exchange applies. This gap may be the result of nonlocal condensates $^{34,40,41}$, which have a finite fluctuation size, or alternatively being induced in the form of an effective quark mass acquired through the interaction with an instanton background $^{42,43}$. But the general result is the same: the shape of the pion distribution amplitude gets strongly enhanced in the central region and resembles closely the asymptotic one.

In view of these drawbacks, a potentially good agreement between theoretical estimates employing CZ-type distribution amplitudes for the pion – as recently reported in $^{28}$ – and experimental measurements is entirely circumstantial.
In the present work, we are going to show that taking together the soft form-factor contribution due to the overlap of the initial and final pion wave functions \[24\], and the hard, i.e., factorizing part, we can obtain a result in qualitative agreement with the existing data and complying with the power counting rules. Actually, including also NLO contributions to the hard-scattering amplitude (see, \[44\] and previous references therein), \(F_\pi(Q^2)\) gets additionally enhanced to account for approximately half of the form-factor magnitude \[11\], modulo the large uncertainties of the existing data.

However, we cannot and do not exclude that the true pion distribution amplitude may deviate from the asymptotic one, but this deviation should be within the margins allowed by the experimental data for the \(\pi \gamma\) transition form factor. Hence the true pion distribution amplitude may well be a hybrid of the type \(\Phi_{\text{true}} = 90\% \Phi_{\text{as}} + 9\% \Phi_{\text{CZ}} + 1\% C_3^4\). This mixing ensures a broader shape of the pion distribution amplitude, with the fourth-order, “Mexican hat”-like, Gegenbauer polynomial, being included in order to cancel the dip at \(x = 1/2\). The shapes derived from instanton-based approaches \[24,43\] are of this type. For such distribution amplitudes, evolution already at LO must be taken into account that tends to reduce the importance of the endpoint region leading to a decrease of the magnitude of the form factors towards the data \[24,44\]. On the other hand, for the asymptotic solution, evolution enters only at NLO and is a tiny effect which is ignored in the present exploratory investigation.

The starting point of our analysis is the expression for the pion form factor in the transverse configuration space after employing factorization to separate a short-distance, i.e., hard-scattering part (where the terminology of \[15,13,14\] is adopted):

\[
F_\pi(Q^2) = \int_0^1 \! \! \! dx \, dy \int_{-\infty}^{\infty} \frac{d^2 b}{(4\pi)^2} P_{\text{out}}^\pi(y, b, P', \mu_F, \mu_R) T_H(x, y, b, Q/\mu_R, Q/\mu_F) \times P_{\text{in}}^\pi(x, b, P, \mu_F, \mu_R),
\]  

(2)

Here \(P^+ = Q/\sqrt{2} = P^- \), \(Q^2 = - (P' - P)^2\), and \(\mu_R = C_2 \xi Q\) and \(\mu_F = C_1/b\) are, respectively, the renormalization and factorization scales, with \(C_1, C_2 \sqrt{2} = C_2^{CS} \) \[15\] (\(\xi = x, \bar{x}, y, \bar{y}\)), the constants \(C_1, C_2\) being integration constants of order unity, so that (uncalculated) higher-order corrections are small \[15,18\]. Finally, \(b\) is the variable conjugate to the transverse gluon momentum, and denotes the transverse distance between quark and antiquark.

The hard-scattering amplitude \(T_H\) is the amplitude for a quark and an antiquark to scatter collinearly via a hard-gluon exchange with wavelengths limited by \(b\), and is given in leading order by

\[
T_H(x, y, b, Q/\mu_R, bQ) = 8C_F a_s^{an}(\mu_R^2) K_0(\sqrt{x y b Q}).
\]  

(3)

In Eq. (2), \(P_{\pi}\) describes the valence \(q \bar{q}\) amplitude which includes gluonic radiative corrections \[14\] as well as the primordial transverse size of the bound state \[24\]:

\[
P_{\pi}(x, b, Q, \mu_F, \mu_R) = \exp \left[ -s(x, b, Q, \mu_F, \mu_R) - s(\bar{x}, Q, \mu_F, \mu_R) - 2 \int_{\mu_F}^{\mu_R} \frac{d\mu}{\mu} \gamma_q(g(\mu)) \right] \times P(x, b, \mu_F).
\]  

(4)

The pion distribution amplitude at the factorization point is approximately given by
\[ \mathcal{P} (x, b, \mu_F = C_1/b) \simeq \phi_\pi (x, \mu_F = C_1/b) \Sigma(x, b) , \]

where \( \Sigma(x, b) \) parametrizes the intrinsic transverse size of the pion (see below). In the collinear approximation, one has

\[ \frac{f_\pi}{2\sqrt{2N_c}} \phi_\pi (x, \mu_F^2) = \int_{\mu_R}^{\mu_F} \frac{d^2k_\perp}{16\pi^3} \Psi_\pi (x, k_\perp) , \]

where \( \mu_F = 130.7 \text{ MeV} \) and \( N_c = 3 \). Integrating on both sides of this equation over \( x \) normalizes \( \phi_\pi \) to unity, i.e., \( \int_0^1 dx \phi_\pi (x, \mu_F^2) = 1 \) because the rhs is fixed to \( \frac{f_\pi}{2\sqrt{2N_c}} \) by the leptonic decay \( \pi \rightarrow \mu^+ \nu_\mu \) for any factorization scale.

The Sudakov functions can be written in terms of the momentum-dependent cusp anomalous dimension to read \([17,18,19,20]\)

\[ s (\xi, b, Q, \mu_F, \mu_R) = \frac{1}{2} \int_{\mu_F}^{\mu_R=\xi} \frac{d\mu}{\mu} \Gamma_{cusp} (\gamma, g(\mu)) , \]

where \( \gamma = \ln \left( \frac{C_2Q}{\mu} \right) \) is the cusp angle, i.e., the emission angle of a soft gluon and the bent quark line after injecting at the cusp point the external (large) momentum by the off-mass-shell photon, and

\[ \Gamma_{cusp} (\gamma, g(\mu)) = 2 \ln \left( \frac{C_2\xi Q}{\mu} \right) A (g(\mu)) + B (g(\mu)) . \]

The functions \( A \) and \( B \) are defined by

\[ A (g(\mu)) = \frac{1}{2} \left[ 2\Gamma_{cusp} (g(\mu)) + \beta(\mu) \frac{1}{\partial g} \mathcal{K}(C_1, g(\mu)) \right] \]

\[ = C_F \frac{\alpha_s^a(g(\mu))}{\pi} + \frac{1}{2} K C_F \left( \frac{\alpha_s^a(g(\mu))}{\pi} \right)^2 , \]

and

\[ B (g(\mu)) = -\frac{1}{2} \left[ \mathcal{K}(C_1, g(\mu)) + \mathcal{G}(\xi, C_2, g(\mu)) \right] \]

\[ = \frac{2}{3} \frac{\alpha_s^a(g(\mu))}{\pi} \ln \left( \frac{C_F^2}{C_2^2} \frac{e^{2\gamma_E-1}}{4} \right) , \]

respectively, and the \( K \)-factor in the \( \overline{\text{MS}} \) scheme is given by the expression

\[ K = \left( \frac{67}{18} - \frac{\pi^2}{6} \right) C_A - \frac{10}{9} n_f T_F + \frac{1}{8} \beta_0 \ln \left( C_1 e^{\gamma_E}/2 \right) \]

where \( C_A = N_c = 3, n_f = 3, T_F = 1/2, \gamma_E \) being the Euler-Mascheroni constant. The functions \( \mathcal{K}, \mathcal{G} \) are calculable within perturbative QCD \([15]\). Both anomalous dimensions, \( \Gamma_{cusp}(g(\mu)) = C_F \frac{\alpha_s^a(\mu^2)}{\pi} \) and \( \gamma_q(g(\mu)) = -\frac{\alpha_s^a(\mu^2)}{\pi}, \) will be evaluated using the analytic model of \([1]\) in next-to-leading logarithmic order (for more technical details, we refer to \([11]\)).
the long-dashed line which shows the soft contribution \cite{24}, in comparison with the existing experimental spacelike pion form factor calculated with $F_\pi$ within different factorization schemes, except for the long-dashed line which shows the soft contribution \cite{24}, in comparison with the existing experimental data \cite{22,23}. The lower solid line shows the IR-enhanced result obtained in our analysis, whereas the upper one stands for the sum of the long-dashed and the lower solid line. The dot-dashed line reproduces the calculation of \cite{26}, which does not include an intrinsic $k_\perp$-dependence.
For simplicity, we follow [21] and model the distribution of primordial (intrinsic) transverse momentum in the pion wave function in the form of a Gaussian normalized to unity:

$$\Psi_\pi(x, k_\perp) = \frac{16\pi^2 f_\pi}{2\sqrt{2N_c}} \phi(x) \beta^2 g(x) \exp \left[ -g(x) \beta^2 k_\perp^2 \right],$$

where $g(x) = 1/x\bar{x}$ and the quark masses are neglected. For $\phi_{\text{av}}(x) = 6x\bar{x}$, one has $\beta^2 = 0.883 \text{[GeV}^{-2}]$ which corresponds to $<k_\perp^2>^{1/2} = 350 \text{ MeV}$.

Before we proceed with the presentation of our results, exposed in Fig. [1], let us at this point interject some comments regarding the role of the scales entering the calculation of the pion form factor. Whenever $\xi < \frac{\xi}{C_2^z Q}$, all Sudakov exponential factors are set equal to unity [14]. For all values of $b$, there is a hierarchy of scales according to $\Lambda_{\text{QCD}}^z \ll \mu_F = C_1/b \leq \mu_R = C_2\xi Q \leq Q$. The limit $\mu_R \simeq \mu_F$ can be interpreted as the minimum virtuality scale of exchanged quanta (or equivalently as the maximum transverse separation) in $T_\text{H}$ below which propagators cannot be treated within perturbation theory and are therefore absorbed into $\phi_\pi$. In the present analysis we use $\mu_F = C_1/b$, $\mu_R = C_2\xi Q$ with $\xi = x, \bar{x}, y, \bar{y}$, and $C_1 = \exp \left[ -\frac{1}{2}(2\gamma_E - 1) \right]$, $C_2 = 1/\sqrt{2}$, the latter corresponding to the value $C_2^z = 1$ (cf. [13]). For this choice of the scheme constants, the logarithmic term in the $K$ factor is eliminated. Note that for technical reasons, the argument of $\alpha_s$ in $T_\text{H}$ is taken to be $\mu_R = C_2\sqrt{\xi}Q$, where $\xi$ is either $x$ or $y$.

In order to show how the contribution to the pion form factor is accumulated in $b$ space, we show in Fig. [2] the dependence of the scaled pion form factor against $b\Lambda_{\text{QCD}}$. One observes a fast rise of the displayed curves as $Q$ increases. Indeed, already for the smallest value shown, $Q = 2 \text{ GeV}$, the form factor accumulates half of the whole contribution in the region $b \simeq 0.5/\Lambda_{\text{QCD}}$. For still larger $Q$ values, the form factor levels off already around $b = 0.3/\Lambda_{\text{QCD}}$ for $Q = 5 \text{ GeV}$, and $b = 0.25/\Lambda_{\text{QCD}}$ for $Q = 10 \text{ GeV}$. This behavior of the curves uncovers how the IR stability and self-consistency of the perturbative treatment in the present scheme is improved, as compared to previous, conventional, approaches [13][21][28].

A similar expression to Eq. (2) holds also for $F_{\pi^0 \gamma^* \gamma}$, the main difference being that the latter contains only one pion wave function, and furthermore the associated short-distance part, $T_\text{H}$, does not depend on $\alpha_s$ in LO. The only dependence on the strong coupling constant at leading order enters through the anomalous dimensions of the cusp and the quark wave function. The result of this calculation is displayed in Fig. [3]. Notice that in this case, we use $\mu_R = C_2\sqrt{\xi}Q$, where $\xi = x$ or $\xi = \bar{x}$.

In summary, we have shown that modifying $\alpha_s$ in the IR region by a nonperturbative power correction, which removes the unphysical Landau singularity, may play a crucial role in the practical calculation of exclusive processes because it improves the IR stability of computed observables based on perturbation expansions without introducing external parameters to “freeze” the running strong coupling. Furthermore, we have given quantitative evidence that in this way it is possible to get an enhanced hard contribution to $F_\pi(Q^2)$, relying exclusively on the asymptotic form of the pion distribution amplitude, so that, though this contribution comprises Sudakov corrections and a primordial $k_\perp$-dependence, it is not suppressed. Together with the soft part of $F_\pi(Q^2)$, this contribution can account for the trend of the existing (admittedly low-accuracy) data without employing endpoint-concentrated pion distribution amplitudes. The same treatment yields for $F_{\pi^0 \gamma^* \gamma}$ a theoretical prediction which is in good agreement with the data.
Fig. 2. Dependence on the (critical) transverse distance $b$ of the hard contribution to the scaled pion form factor, calculated with $\phi_{as}$ within our IR-finite factorization scheme, for three different values of the momentum transfer: $Q_1 = 2$ GeV (solid line), $Q_2 = 5$ GeV (dashed line), and $Q_3 = 10$ GeV (dashed-dotted line).
Fig. 3. Pion-photon transition form factor calculated with $\phi_{as}$ and IR enhancement (lower solid line). The other solid line shows the asymptotic behavior. The data are taken from $^{25,26}$.

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