Direct photons and dileptons via color dipoles

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I. INTRODUCTION

Massive lepton pair production and inclusive direct photon production in hadronic collisions have historically provided an important tool to gain access to parton distributions in hadrons. Moreover, direct photons, i.e. photons not from hadronic decay, can be also a powerful probe of the initial state of matter created in heavy ion collisions, since they interact with the medium only electromagnetically and therefore provide a baseline for the interpretation of jet-quenching models.

In the parton model, the Feynman diagrams for partonic subprocesses that are present in Drell-Yan (DY) lepton pair production and in inclusive direct photon production are different, and the connection between both production mechanisms within a unique approximation scheme is not obvious. Since in the target rest frame the DY process looks like bremsstrahlung of a virtual photon decaying into a lepton pair, we will show that the color dipole formalism defined in this frame is well suited to describe both production processes in a unified framework free of parameters. As an illustrative example, we study dilepton spectra in 800-GeV pp collisions, inclusive direct photon spectra for pp collisions at RHIC energies \( \sqrt{s} = 200 \text{ GeV} \), and for \( p\bar{p} \) collisions at Tevatron energies \( \sqrt{s} = 1.8 \text{ TeV} \), in a formalism that is free from any extra parameters.

This letter is an alternative attempt. We shows that the color dipole approach can successfully describe inclusive photon production in hadron-hadron collisions.

II. COLOR DIPOLE FORMALISM

The color dipole formalism, developed in [3] for the case of the total and diffractive cross sections, can be also applied to radiation [3]. Although in the process of electromagnetic bremsstrahlung by a quark no dipole participates, the cross section can be expressed via the more elementary cross section \( \sigma_{qq} \) of interaction of a \( q\bar{q} \) dipole. Nevertheless, this is a fake, or effective dipole. Similar to a real dipole, where color screening is provided by interactions with either the quark or the antiquark, in the case of radiation the two amplitudes for radiation prior or after the interaction screen each other, leading to cancellation of the infra-red divergences [7].

The transverse momentum \( p_T \) distribution of photon bremsstrahlung in quark-nucleon interactions, integrated over the final quark transverse momentum, was derived in [8] in terms of the dipole formalism,

\[
\frac{d\sigma^N(q \rightarrow q\gamma)}{d(ln\alpha)d^2\vec{p}_T} = \frac{1}{(2\pi)^2} \sum_{i,n,f,L,T} \int d^2\vec{r}_1 d^2\vec{r}_2 e^{i\vec{p}_T.(\vec{r}_1 - \vec{r}_2)} \times \phi^{T,L}_q(\alpha, \vec{r}_1)\phi^{T,L}_{\bar{q}}(\alpha, \vec{r}_2) \Sigma_\gamma(x, \vec{r}_1, \vec{r}_2, \alpha),
\]

where

\[
\Sigma_\gamma(x, \vec{r}_1, \vec{r}_2, \alpha) = \frac{1}{2} \{ \sigma_{qq}(x, \alpha r_1) + \sigma_{qq}(x, \alpha r_2) \}
- \frac{1}{2} \sigma_{q\bar{q}}(x, \alpha (\vec{r}_1 - \vec{r}_2)).
\]

and \( \vec{r}_1 \) and \( \vec{r}_2 \) are the quark-photon transverse separations in the two radiation amplitudes contributing to the cross section, Eq. (1), which correspondingly contains double-Fourier transformations. The parameter \( \alpha \) is the...
relative fraction of the quark momentum carried by the photon, and is the same in both amplitudes, since the interaction does not change the sharing of longitudinal momentum. The transverse displacement between the initial and final quarks is $\alpha r_1$ and $\alpha r_2$ respectively. Since the amplitude of quark interaction has a phase factor $\exp(i\vec{b} \cdot \vec{p}_T)$, where $\vec{b}$ is the impact parameter of collision, the transverse displacement between the initial and final quarks leads to the color screening factor $1 - \exp(i\alpha r^2 \vec{p}_T)$. In Eq. (1) $T$ stands for transverse and $L$ for longitudinal photons. The energy dependence of the dipole cross section which comes via the variable $\e$. The energy dependence of the dipole cross section, when $\alpha r_1 \rightarrow \infty$, is the projectile four-momentum and $q$ is the four-momentum of the dilepton, is generated by additional radiation of gluons which can be resummed in the leading $\ln(1/x)$ approximation.

In Eq. (1) the light-cone (LC) wavefunction of the projectile quark $\gamma q$ fluctuation has been decomposed into transverse $\phi_{\gamma q}^T(\alpha, \vec{r})$ and longitudinal $\phi_{\gamma q}^L(\alpha, \vec{r})$ components, and an average over the initial quark polarization and sum over all final polarization states of quark and photon is performed. These LC wavefunction components $\phi_{\gamma q}^T(\alpha, \vec{r})$ can be represented at the lowest order as:

$$
\sum_{i,n,f} \phi_{\gamma q}^T(\alpha, \vec{r}_1) \phi_{\gamma q}^T(\alpha, \vec{r}_2) = \frac{\alpha_{em}}{2\pi^2} m_q^2 \alpha^4 K_0(\alpha) K_0(\alpha)
$$

$$
+ \frac{\alpha_{em}}{2\pi^2} (1 + (1-\alpha)^2) \frac{\vec{r}_1 \cdot \vec{r}_2}{r_1 r_2} K_1(\alpha) K_1(\alpha),
$$

$$
\sum_{i,n,f} \phi_{\gamma q}^L(\alpha, \vec{r}_1) \phi_{\gamma q}^L(\alpha, \vec{r}_2) = \frac{\alpha_{em}}{\pi^2} M^2 (1 - \alpha)^2
$$

$$
\times K_0(\alpha) K_0(\alpha),
$$

(3)

in terms of transverse separation $\vec{r}$ between photon $\gamma$ and quark $q$ and the relative fraction $\alpha$ of the quark momentum carried by the photon. Here $K_0(1)$ denotes the modified Bessel function of the second kind. We have also introduced the auxiliary variable $\alpha^2 = \frac{1}{2} \frac{m_q^2}{\alpha^2} + (1 - \alpha) M^2$, where $M$ denotes the mass of dilepton and $m_q$ is an effective quark mass which can be conceived as a cutoff parameter. This quark mass has less influence on dilepton production in $pp$ collisions, albeit it will be a numerically important parameter for direct photon production, when $M = 0$. In general the quark mass $m_q$ should not be considered an extra parameter. Indeed, depending on the kinematical variable $M$, the Feynman variable $x_F$ and the square of the center of mass energy of the colliding hadrons $s$, there always exists a range of values of $m_q$ where the result does not depend on the specific $m_q$ value. For direct photon $M = 0$, $m_q$ cannot be zero since the wave function becomes divergent. In this paper, as in Refs. [3, 4], we take $m_q = 0.2$ GeV for both dilepton and direct photon production. Notice also that $m_q$ is a more important parameter in proton-nucleus collisions where a value of $m_q = 0.2$ GeV is needed in order to describe the nuclear shadowing effect [10].

In order to obtain the hadron cross section from the elementary partonic cross section Eq. (1), one should sum up the contributions from quarks and antiquarks weighted with the corresponding parton distribution functions (PDFs) $[8, 9]$.

$$
\frac{d\sigma^{DY}(pp \rightarrow \gamma^* X)}{dM^2 dx_F d^2p_T} = \frac{\alpha_{em}}{3\pi M^2 x_1 + x_2} \int_{x_1}^{1} \frac{dx}{\alpha^2}
$$

$$
+ \sum_{i,n,f} Z_i^2 f_i^2(\frac{x}{\alpha}) + \bar{q}_i(\frac{x}{\alpha}) \frac{d\sigma^{\gamma N}(q \rightarrow q\gamma^*)}{d(ln\alpha) d^2p_T}
$$

$$
= \frac{\alpha_{em}}{3\pi M^2(x_1 + x_2)} \int_{x_1}^{1} \frac{dx}{\alpha F_{q}^p(\frac{x}{\alpha}), Q} \frac{d\sigma^{\gamma N}(q \rightarrow q\gamma^*)}{d(ln\alpha) d^2p_T}.
$$

(4)

The PDFs of the projectile enter in a combination which can be written in terms of proton structure function $F_2^p$. Notice that with our definitions the fractional quark charge $Z_i$ is not included in the LC wave function of Eq. (3), and that the factor $\frac{\alpha_{em}}{\pi^2}$ in Eq. (4) accounts for the decay of the photon into the lepton pair. We use the standard notation for the kinematical variables, $x_1 = (\sqrt{x_F^2 + 4\tau + x_F})/2$ denotes the momentum fraction that the photon carries away from the projectile hadron in the target frame, we define $x_2 = x_1 - x_F$, $x_F = 2p_L/\sqrt{s}$ is the Feynman variable and $\tau = \frac{M^2 + s}{2p_L}$, where $p_L$ and $p_T$ denote the longitudinal and transverse momentum components of the photon in the hadron-hadron center of mass frame, $s$ is the center of mass energy squared of the colliding protons and $M$ is the dilepton mass. We also need to identify the scale $Q$ entering in the proton structure function in Eq. (5), and relate the energy scale $x$ of the dipole cross section entered in Eq. (2) to measurable variables. From our previous definition, and following previous works [8, 11] we have that $x = x_2$. At zero transverse momentum, the dominant term in the LC wavefunction Eq. (3) is the one that contains the modified Bessel function $K_1(\alpha r)$. This function decay exponentially at large values of the argument, so that the mean distances which numerically contribute are of order $1/\alpha$. On the other hand, the minimal value of $\alpha$ is $x_1$, and therefore the virtuality $Q^2$ which enters into the problem at zero transverse momentum is $\approx (1 - x_1)M^2$. Thus the hard scale at which the projectile parton distribution is probed turns out to be $Q^2 = p_T^2 + (1 - x_1)M^2$. Notice that in the previous studies, $M^2$ [9] and $(1 - x_1)M^2$ [11] were used for the scale $Q^2$. Nevertheless, these different choices for $Q^2$ bring less than about a 20% effect at small $x_2$ values.

The dipole cross section is theoretically unknown, although several parametrizations have been proposed in the literature. For our purposes, here we consider two parametrizations, the saturation model of Golec-Biernat and Wüsthoff (GBW) [12] and the modified GBW coupled to DGLAP evolution (GBW-DGLAP) [13].
A. GBW model parametrization

In the GBW model \[12\] the dipole cross section is parametrized as,

\[
\sigma_{q\bar{q}}(x, r) = \sigma_0 \left( 1 - e^{-r^2/R_0^2} \right),
\]

where the parameters, fitted to DIS HERA data at small \(x\), are given by \(\sigma_0 = 23.03\) mb, \(R_0 = 0.4\)fm \(\times (x/x_0)^{0.144}\), where \(x_0 = 3.04 \times 10^{-4}\). This parametrization gives a quite good description of DIS data at \(x < 0.01\). A salient feature of the model is that for decreasing \(x\), the dipole cross section saturates for smaller dipole sizes, and that at small \(r\), as perturbative QCD implies, \(\sigma \sim r^2 \) vanishes. This is the so-called color transparency phenomenon \[3\].

One of the obvious shortcomings of the GBW model is that it does not match with QCD evolution (DGLAP) at large values of \(Q^2\). This failure can be clearly seen in the energy dependence of \(\sigma_{\text{tot}}^{q\bar{q}}\) for \(Q^2 > 20\) GeV\(^2\), where the the model predictions are below the data \[12,13\].

B. GBW couple to DGLAP equation and dipole evolution

A modification of the GBW dipole parametrization model, Eq. (6), was proposed in Ref. \[13\]

\[
\sigma_{q\bar{q}}(x, r) = \sigma_0 \left( 1 - \exp \left( -\frac{\pi^2 r^2 \alpha_s(x)}{3\sigma_0} \right) \right),
\]

where the scale \(\mu^2\) is related to the dipole size by

\[
\mu^2 = \frac{C}{r^2} + \mu_0^2.
\]

Here the gluon density \(g(x, \mu^2)\) is evolved to the scale \(\mu^2\) with the leading order (LO) DGLAP equation \[14\]. Moreover, the quark contribution to the gluon density is neglected in the small \(x\) limit, and therefore

\[
\frac{\partial g(x, \mu^2)}{\partial \ln \mu^2} = \frac{\alpha_s(\mu^2)}{2\pi} \int_x^1 dz P_{gg}(z) \frac{x}{z} g(x, \mu^2) - g(z, \mu^2).
\]

where \(P_{gg}(z)\) and \(\alpha_s(\mu^2)\) denote the QCD splitting function and coupling, respectively. The initial gluon density is taken at the scale \(Q_0^2 = 1\) GeV\(^2\) in the form

\[
x g(x, \mu^2) = A_g x^{-\lambda_g}(1-x)^{5.6},
\]

where the parameters \(C = 0.26, \mu_0^2 = 0.52\) GeV\(^2\), \(A_g = 1.20\) and \(\lambda_g = 0.28\) are fixed from a fit to the DIS data for \(x < 0.01\) and in a range of \(Q^2\) between 0.1 and 500 GeV\(^2\) \[13\]. We use the LO formula for the running of the strong coupling \(\alpha_s\), with three flavors and for \(A_{\text{QCD}} = 0.2\) GeV. The dipole size determines the evolution scale \(\mu^2\) through Eq. (8). The evolution of the gluon density is performed numerically for every dipole size \(r\) during the integration of Eq. (1). Therefore, the DGLAP equation is now coupled to our master equations (15). It is important to stress that the GBW-DGLAP model preserves the successes of the GBW model at low \(Q^2\) and its saturation property for large dipole sizes, while incorporating the evolution of the gluon density by modifying the small-\(r\) behaviour of the dipole size.

The proton structure function in Eqs. (4,5) is parametrized as

\[
F_2^p(x, Q^2) = A(x) \left[ \frac{\ln(Q^2/\Lambda^2)}{\ln(Q_0^2/\Lambda^2)} \right] B(x) \left( 1 + \frac{C(x)}{Q^2} \right),
\]

with \(Q_0^2 = 20\) GeV\(^2\), \(\Lambda = 0.25\) GeV, and the functions \(A(x), B(x)\) and \(C(x)\) are parametrized in terms of 17 parameters fitted to different experiments, and whose functional forms can be found in the Appendix of Ref. \[13\]. This parametrization is only valid in the kinematic range of the data sets which cover correlated regions in the ranges \(3.5 \times 10^{-5} < x < 0.85\) and \(0.2 < Q^2 < 5000\) GeV\(^2\).

III. NUMERICAL RESULTS

Before we proceed to present the results in the color dipole approach, some words regarding the validity of this formulation are in order. Although both valence and sea quarks in the projectile are taken into account through the proton structure function Eqs. (4,5), the color dipole picture accounts only for Pomeron exchange from the target, while ignoring its valence content. In terms of Regge phenomenology, this means that Reggeons are not taken into account, and as a consequence, the dipole approach predicts the same cross sections for both particle and antiparticle induced DY reactions. Therefore, in principle this approach is well suited for high-energy processes, i.e. small \(x_2\). The exact range of validity of the dipole approach is of course not known a priori, but there is evidence \[2,11\] in its favor for values of \(x_2 < 0.1\). In our case, however, we use a parametrization of the dipole cross section fitted to DIS data for Bjorken-\(x < 0.01\) and for energy scales \(Q^2 < 500\). Given these restrictions, at present there are not many data for DY cross section at low \(x_2\). Notice also that some data are integrated over \(x_F\) and \(M\), and are therefore contaminated by contributions not included into the color dipole approach.

We compare the present approach to data for 800-GeV pp collisions from E886 \[1\], which are not integrated over \(x_F\) and \(M\), and correspond to the lowest \(x_2\) values, i.e. lightest \(M\) and highest \(x_F\). We selected a \(x_F\) bin where \(0.55 < x_F < 0.8\), with an average value \(\langle x_F \rangle = 0.63\). Within this bin we selected two bins with the lightest average values for \(M\), one for \(4.20 < M_{\mu^+\mu^-} < 5.20\), with an average value of \(\langle M_{\mu^+\mu^-} \rangle = 4.80\) GeV, and the other for \(5.20 < M_{\mu^+\mu^-} < 6.20\), with an average value \(\langle M_{\mu^+\mu^-} \rangle = 5.70\) GeV. The experimental data are plotted, with errors, in Figs. (1) and (2). In our calculations we have taken the experimental average values for \(x_F\) and \(M\).
FIG. 1: The Dilepton spectrum in 800-GeV $pp$ collisions at $x_F = 0.63$ and $M = 5.7$ GeV. We show the result of the GBW dipole model (dashed line) and the GBW-DGLAP model (dotted line). We also show the result when a constant primordial momentum $\langle k_T^2 \rangle = 0.4$ GeV$^2$ is incorporated within the GBW-DGLAP dipole model (solid line). Experimental data are from Ref. [1]. The E866 error bars are the linear sum of the statistical and systematic uncertainties. An additional $\pm 6.5\%$ uncertainty in the experimental data points due to the normalization is also common to all points.

In Figs. 1, 2 we show the result obtained by the dipole approach, for both the GBW and the GBW-DGLAP dipole models. At low transverse momentum $p_T < 2$ GeV both model predictions are almost identical, but at higher $p_T$ the dipole parametrization improved by DGLAP evolution bends down towards the experimental points improving the result. This is more obvious for higher values of $M$. Notice that for the case of a smaller value of $x_F$ with a lighter $M$, Fig. 2 where the dipole approach is better suited, the GBW model without inclusion of the DGLAP evolution already provides a good description of the data. We stress that the theoretical curves in Figs. 1, 2 are the results of a parameter free calculation. As we already pointed out, varying the quark mass $m_q$ leaves the numerical results almost unaffected. Notice also that in contrast to the LO parton model, no $K$-factor was introduced, since the dipole parametrization fitted to DIS data already includes contributions from higher order perturbative corrections as well as non-perturbative effects contained in DIS data.

One of the data points which surprisingly is left out from our theoretical computation curves for both values of $M$, is the one at the lowest $p_T$. In the dipole approach the DY cross section is finite at $p_T = 0$ due to the saturation of the dipole cross section, which is in striking contrast to the LO pQCD correction to the parton model, where one needs to resume the large logarithms $\ln(p_T^2/M^2)$ from soft gluon radiation in order to obtain a physically sensible results at $p_T = 0$ [16]. One of the possible reasons behind the lack of agreement between our result and the experimental data at $p \to 0$ may be due to a soft non-perturbative primordial transverse momentum distribution of the partons in the colliding protons. Such a primordial transverse momentum may have various non-perturbative origins, e. g. finite size effects of the hadron, instanton effects, pion-cloud contributions. Moreover, in the parton model it has been shown that even within the next-to-leading order (NLO) pQCD correction, experimental data of heavy quark pair production [17], direct photon production [18] and DY lepton pair production [19] can be only described if an average primordial momentum as large as 1 GeV is included (see also Ref. [20]). Such a large value for the initial transverse momentum strongly indicates its perturbative origin in the parton model, and in principle must have been already included in the pQCD correction. Therefore, in the pQCD approach, it is still an open question how to separate what is truly intrinsic and what is pQCD generated transverse momentum. However, in the dipole approach all perturbative and non-perturbative contributions, apart from the finite-size effect of hadrons, are already encoded into the cross section via fitting the dipole parameters to DIS data. Therefore, we expect that in the dipole approach the primordial momentum should have a purely non-perturbative origin, and to be considerably less than in the parton model. One may introduce an intrinsic momentum contribution in the following factorized form

$$\mathcal{F}(p_T) \rightarrow \int d^2k_T \mathcal{F}(p_T - k_T) G_N(k_T), \quad (12)$$
where the function $F$ denotes the cross section defined in Eqs. (15). We assume that the initial $p_T$ distribution $G_N(p_T)$ has a Gaussian form,

$$G_N(k_T) = \frac{1}{\pi(k_T^2/k_T^2)_N} e^{k_T^2/(k_T^2)_N},$$

(13)

where $\langle k_T^2 \rangle_N$ is the square of the two-dimensional width of the $p_T$-distribution for an incoming quark, and also that $\langle k_T^2 \rangle_N$ is a constant independent of the hard scale $Q$, since the pQCD radiation-generating transverse momenta are already taken into account in our approach. The differential cross section convoluted with the primordial momentum distribution in the GBW-DGLAP dipole model are shown in Figs. 1 and 2 with curves denoted with GBW-DGLAP-Primordial. A value around $\langle k_T^2 \rangle_N = 0.4$ GeV$^2$ can describe the experimental points at low $p_T$ for both sets of data plotted in Figs. 1 and 2. This value, as we expected, is lower than the primordial momentum which has been used in the parton model.

The experimental data points for $p_T \to 0$ should be taken with some precaution, since there exists some disagreement between different experiments for DY lepton pair production at low $p_T$. Indeed, although the E772 and E866 measurements [1] have good agreement among them over a wide range of values, they disagree at $p_T \to 0$. Therefore, the discrepancy between our theoretical results and experimental data at $p_T \to 0$ might be just in fact an artifact of the experiments.

Next we calculate the inclusive direct photon spectra within the same framework. For direct photon we have $M = 0$, and we assume again a quark mass $m_q = 0.2$ GeV. As illustrative examples we compare our results with the PHENIX and CDF experiments. Notice that direct photon problem (with $M = 0$), compared to the massive virtual photon case (with $M$ as big as $\sim 5$ GeV), is numerically more involved since the integrand in Eq. (11) is divergent when $m_q \to 0$.

In Fig. 3 we show the differential cross section obtained from the GBW and the GBW-DGLAP dipole models at midrapidities, for $pp$ collisions at RICH energies $\sqrt{s} = 200$ GeV. The experimental data are from the PHENIX measurements for inclusive direct photon production at $y = 0$ [2]. We have also checked out that the effect of the incorporation of the same transverse primordial momentum $\langle k_T^2 \rangle_N = 0.4$ GeV$^2$ which can describe the dilepton spectra at low $p_T$, will be in this case too small to improve the results at the range of $p_T$ of the experimental data. Without a physically sound guiding principle, however, the introduction of a higher value of intrinsic momentum is somehow unsatisfactory and will not be further discussed here. Notice also that, in contrast to the parton model, we have not included any photon fragmentation function [21, 22, 23] for computing the cross section, since the dipole formulation already incorporates all perturbative (via Pomeron exchange) and non-perturbative radiation contributions. It has been shown that the NLO pQCD prediction [21, 23] are also consistent with the RHIC data within the uncertainties [2].

In Fig. 4 we show the dipole models predictions for inclusive prompt-photon production at midrapidities $|y| = 0$ at RHIC energy $\sqrt{s} = 200$ GeV. Experimental data are from Ref. [2]. In the down panel we use the GBW-DGLAP dipole model result for the theory. The error bars are the linear sum of the statistical and systematic uncertainties.
CDF energy and GBW and GBW-DGLAP dipole models for midrapidity at CDF energy $\sqrt{s} = 1.8$ TeV. Experimental data are for inclusive isolated photon from CDF experiment for $p\bar{p}$ collision at CDF energy and $|\eta| < 0.9$ \cite{2}. The NLO QCD curve is from the authors of reference \cite{26} (given in table 3 of Ref. \cite{26}) and use the CTEQ5M parton distribution functions with the all scales set to the $p_T$. In the down panel we use the GBW-DGLAP dipole model result for the theory. The error bars are the linear sum of the statistical and systematic uncertainties.

FIG. 4: Inclusive direct photon spectra obtained from the CDF experiment for the color dipole approach without any free parameters. In contrast to the parton model, in the dipole approach there is no ambiguity in defining the intrinsic transverse momentum. Such a purely non-perturbative primordial momentum improves the results in the case of dilepton pair production, but does not play a significant role for direct photon production at the given experimental range of $p_T$. We also showed that the color dipole formulation coupled to the DGLAP evolution provides a better description of data at large transverse momentum compared to the GBW dipole model.

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