Phenomenology of TMDs Using SCET

Stefano Melis*

*Dipartimento di Fisica, Università di Torino, Via Pietro Giuria 1,
Torino, I-10125, Italy
stefano.melis@to.infn.it*

Umberto D'Alesio

*Dipartimento di Fisica, Università di Cagliari, Italy
INFN, Sezione di Cagliari, C. P. 170
Monserrato (CA), I-09042, Italy
umberto.dalesio@ca.infn.it*

Miguel G. Echevarria

*Nikhef and Department of Physics and Astronomy,
VU University Amsterdam, De Boelelaan 1081,
HV Amsterdam, NL-1081, the Netherlands
m.g.echevarria@nikhef.nl*

Ignazio Scimemi

*Departamento de Física Teórica II,
Universidad Complutense de Madrid (UCM),
Avenida Complutense s/n, Madrid, E-28040, Spain
ignazios@fis.ucm.es*

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We present a global fit of current available data for Drell-Yan and Z -boson production using the Transverse Momentum Dependent evolution formalism of Refs. [1–4]. We performed our analysis including perturbative calculable contributions up to next-to-next-to-leading-logarithmic accuracy. We took into account non-perturbative corrections both to the intrinsic nucleon structure and to the evolution kernel.

Keywords: TMD evolution; resummation; Drell-Yan.

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*Speaker.

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1. Introduction and formalism

In this contribution we present a study of Drell-Yan (DY) and Z -boson production data in the context of the Transverse Momentum Dependent (TMD) evolution scheme developed in Refs. [1–4]. DY and Z -production processes are particularly useful in order to study TMD evolution. In fact, in these processes only TMD Parton Distribution Functions (TMDPDFs) are involved, contrary to other processes like the Semi-Inclusive Deep Inelastic Scattering (SIDIS) where the TMDPDFs appear convoluted with the TMD fragmentation functions. Moreover, present DY/ Z data span a wide region in the mass of the virtual boson, Q , and in its transverse momentum, q_T , thus, allowing the study of the TMDPDFs in both perturbative and non-perturbative regimes.

By means of the factorization theorem, we can write the DY cross section as:

$$\frac{d\sigma}{\sigma_0 dQ^2 dq_T^2 dy} = H(Q^2, \mu^2) \sum_q e_q^2 \int \frac{d^2 \mathbf{b}_T}{4\pi} e^{-i\mathbf{q}_T \cdot \mathbf{b}_T} \times \left\{ \tilde{F}_{q/A}(x_A, b_T; \zeta_A, \mu) \tilde{F}_{\bar{q}/B}(x_B, b_T; \zeta_B, \mu) \right\} + Y, \quad (1)$$

where:

$$\sigma_0 = \frac{4\pi\alpha_{\text{em}}^2}{9Q^2 s}, \quad x_{A,B} = \sqrt{\tau} e^{\pm y}, \quad \tau = \frac{Q^2 + q_T^2}{s} \simeq \frac{Q^2}{s}. \quad (2)$$

Here, $\tilde{F}_{q/N}(x, b_T; \zeta, \mu)$ is the TMDPDF in the b_T space for a parton q in a hadron N ; μ is the factorization scale while ζ is a scale related to the the rapidity divergences that appear in the definition of the TMD soft functions, see Ref. [1] for more details; $H(Q^2, \mu^2)$ is a process dependent hard factor; Y is the part of the collinear cross section (computed at fixed order in α_s) which is regular at $q_T \rightarrow 0$.

Since the TMDPDF contains a non-perturbative part which is relevant at low energy and that must be modeled, it is useful to evolve the TMDPDF from a low initial scale (μ_i, ζ_i), where we can easily set a model, up to an higher scale (μ_f, ζ_f). This can be done by means of an evolver. In impact parameter space we have that

$$\tilde{F}(x, b_T; \zeta_f, \mu_f) = \tilde{R}(b_T; \zeta_i, \mu_i, \zeta_f, \mu_f) \tilde{F}(x, b_T; \zeta_i, \mu_i), \quad (3)$$

where $\tilde{R}$ is the evolver. The expression for the evolver is given by:

$$\tilde{R}(b_T; \zeta_i, \mu_i, \zeta_f, \mu_f) = \exp \left\{ \int_{\mu_i}^{\mu_f} \frac{d\bar{\mu}}{\bar{\mu}} \gamma_F \left(\alpha_s(\bar{\mu}), \ln \frac{\zeta_f}{\bar{\mu}^2} \right) \right\} \left(\frac{\zeta_f}{\zeta_i} \right)^{-D(b_T; \mu_i)}, \quad (4)$$

where γ_F is the anomalous dimension of the TMDPDF and is related to the anomalous dimension of the hard factor as explained in Ref. [1]. The function D follows the renormalization group equation

$$\frac{dD}{d \ln \mu} = \Gamma_{\text{cusp}}, \quad (5)$$

where Γ_{cusp} is the cusp anomalous dimension. Solving this equation with respect to $\mu_b = 2e^{-\gamma_E}/b_T$ one gets:

$$D(b_T; \mu) = D(b_T; \mu_b) + \int_{\alpha_s(\mu_b)}^{\alpha_s(\mu)} d\alpha' \frac{\Gamma_{\text{cusp}}(\alpha')}{\beta(\alpha')}, \quad (6)$$

where $\beta(\alpha_s)$ is the QCD β -function. Writing $\alpha_s(\mu_b)$ as a series in $\alpha_s(\mu)$, one can resum³ all the logs of the type $L_T = \ln(\mu^2/\mu_b^2)$, obtaining a resummed D^R function:

$$\begin{aligned} D^R(b_T; \mu) = & -\frac{\Gamma_0}{2\beta_0} \ln(1-X) \\ & + \frac{1}{2} \left(\frac{a_s}{1-X} \right) \left[-\frac{\beta_1 \Gamma_0}{\beta_0^2} (X + \ln(1-X)) + \frac{\Gamma_1}{\beta_0} X \right] \\ & + \frac{1}{2} \left(\frac{a_s}{1-X} \right)^2 \left[2d_2(0) + \frac{\Gamma_2}{2\beta_0} (X(2-X)) \right. \\ & + \frac{\beta_1 \Gamma_1}{2\beta_0^2} (X(X-2) - 2\ln(1-X)) \\ & \left. + \frac{\beta_2 \Gamma_0}{2\beta_0^2} X^2 + \frac{\beta_1^2 \Gamma_0}{2\beta_0^3} (\ln^2(1-X) - X^2) \right], \end{aligned} \quad (7)$$

where we have used the notation

$$a_s = \frac{\alpha_s(\mu)}{4\pi}, \quad X = a_s \beta_0 L_T, \quad L_T = \ln \frac{\mu^2 b_T^2}{4e^{-2\gamma_E}} = \ln \frac{\mu^2}{\mu_b^2}. \quad (8)$$

This resummation is valid only in a perturbative regime therefore we can write:

$$\tilde{R}^{\text{pert}}(b_T; \zeta_i, \mu_i, \zeta_f, \mu_f) = \exp \left\{ \int_{\mu_i}^{\mu_f} \frac{d\bar{\mu}}{\bar{\mu}} \gamma_F \left(\alpha_s(\bar{\mu}), \ln \frac{\zeta_f}{\bar{\mu}^2} \right) \right\} \left(\frac{\zeta_f}{\zeta_i} \right)^{-D^R(b_T; \mu_i)}. \quad (9)$$

The complete form of the evolution kernel contains a non-perturbative part; we then write:

$$\tilde{R}(b_T; \zeta_i, \mu_i, \zeta_f, \mu_f) = \tilde{R}^{\text{pert}}(b_T; \zeta_i, \mu_i, \zeta_f, \mu_f) \left(\frac{\zeta_f}{\zeta_i} \right)^{-D^{\text{NP}}(b_T)}, \quad (10)$$

where D^{NP} is the non-perturbative piece of the D function. Notice that the relevance of D^{NP} can not be predicted *a priori* and only the experimental data can reveal its role.

In order to make a study of the unpolarized TMDPDF using Eq. (3) we need a model at the initial scale. The main constraint on the structure of TMDPDFs comes from their limit at high transverse momentum, $q_T \gg \Lambda_{\text{QCD}}$. In this case, the

TMDPDFs can be written as the sum of convolutions of Wilson coefficients and collinear Parton Distribution Functions (PDFs):^{1,5}

$$\tilde{F}_{q/N}(x, b_T; \zeta, \mu) = \left(\frac{\zeta}{\mu_b^2} \right)^{-D(b_T; \mu)} \sum_j \int_x^1 \frac{dz}{z} \left\{ \tilde{C}_{q \leftarrow j}^{\mathcal{Q}} \left(\frac{x}{z}, b_T; \mu \right) f_{j/N}(z; \mu) \right\}. \quad (11)$$

The Wilson coefficients $\tilde{C}^{\mathcal{Q}}$ for the quark-quark and the quark-gluon channel at one loop are:¹

$$\begin{aligned} \tilde{C}_{q \leftarrow q'}^{\mathcal{Q}}(z, b_T; \mu) = & \delta(1-z) \delta_{qq'} + 2a_s C_F \left[1 - z - \mathcal{P}_{q \leftarrow q'} L_T \right. \\ & \left. - \delta(1-z) \left(-\frac{1}{2} L_T^2 - \frac{3}{2} L_T + \frac{\pi^2}{12} \right) \right] \delta_{qq'}, \end{aligned} \quad (12)$$

$$\tilde{C}_{q \leftarrow g}^{\mathcal{Q}}(z, b_T; \mu) = 2a_s T_F [2z(1-z) - \mathcal{P}_{q \leftarrow g} L_T], \quad (13)$$

and the one-loop splitting kernels appearing above are given by

$$\mathcal{P}_{q \leftarrow q} = \left(\frac{1+z^2}{1-z} \right)_+, \quad \mathcal{P}_{q \leftarrow g} = z^2 + (1-z)^2. \quad (14)$$

Under renormalization group equation one has:

$$\begin{aligned} \frac{d}{d \ln \mu} \tilde{C}_{q \leftarrow j}^{\mathcal{Q}}(z, b_T; \mu) = & (\Gamma_{\text{cusp}} L_T - \gamma^V) \tilde{C}_{q \leftarrow j}^{\mathcal{Q}}(z, b_T; \mu) \\ & - \sum_i \int_z^1 \frac{d\xi}{\xi} \tilde{C}_{q \leftarrow i}^{\mathcal{Q}}(\xi, b_T; \mu) \mathcal{P}_{i \leftarrow j}(z/\xi). \end{aligned} \quad (15)$$

An explicit expression of γ^V can be found in Ref. [1]. The double logarithms in Eq. (15) can be partially exponentiated^{6,7} obtaining:

$$\tilde{C}_{q \leftarrow j}^{\mathcal{Q}}(z, b_T; \mu) \equiv \exp[h_{\Gamma}(b_T; \mu) - h_{\gamma}(b_T; \mu)] \hat{C}_{q \leftarrow j}(z, b_T; \mu), \quad (16)$$

where h_{Γ} and h_{γ} are defined by the following renormalization equations:

$$\frac{dh_{\Gamma}}{d \ln \mu} = \Gamma_{\text{cusp}} L_T; \quad \frac{dh_{\gamma}}{d \ln \mu} = \gamma^V. \quad (17)$$

Γ_{cusp} , γ^V and $h_{\Gamma, \gamma}$ can be perturbatively expanded in series of α_s , see Refs. [3, 4] for their explicit expressions. We can solve the evolution equation for $h_{\Gamma, \gamma}$ obtaining:

$$\begin{aligned} h_{\Gamma}(b_T; \mu) = & h_{\Gamma}(b_T; \mu_b) + \int_{\mu_b}^{\mu} \frac{d\bar{\mu}}{\bar{\mu}} \Gamma_{\text{cusp}} L_T, \\ h_{\gamma}(b_T; \mu) = & h_{\gamma}(b_T; \mu_b) + \int_{\mu_b}^{\mu} \frac{d\bar{\mu}}{\bar{\mu}} \gamma^V. \end{aligned} \quad (18)$$

The expressions above still contain large logarithms L_T that need to be resummed when $\alpha_s L_T$ is of order 1. We can apply the method of integration used for the

evaluation of D^R to find h_γ^R as the equation for h_γ (second line of Eq. (18)) has the same functional form as Eq. (6), so that

$$\begin{aligned} h_\gamma^R(b_T; \mu) = & -\frac{\gamma_0}{2\beta_0} \ln(1-X) + \frac{1}{2} \left(\frac{a_s}{1-X} \right) \left[-\frac{\beta_1 \gamma_0}{\beta_0^2} (X + \ln(1-X)) + \frac{\gamma_1}{\beta_0} X \right] \\ & + \frac{1}{2} \left(\frac{a_s}{1-X} \right)^2 \left[\frac{\gamma_2}{2\beta_0} (X(2-X)) + \frac{\beta_2 \gamma_0}{2\beta_0^2} X^2 \right. \\ & \quad \left. + \frac{\beta_1 \gamma_1}{2\beta_0^2} (X(X-2) - 2\ln(1-X)) \right. \\ & \quad \left. + \frac{\beta_1^2 \gamma_0}{2\beta_0^3} (\ln^2(1-X) - X^2) \right]. \end{aligned} \quad (19)$$

In the case of h_Γ we note that $L_T = \int_{\alpha_s(\mu_b)}^{\alpha_s(\mu)} \frac{d\alpha}{\beta(\alpha)}$, and in this way Eq. (18) (first line) can be solved as

$$h_\Gamma^R(b_T; \mu) = \int_{\alpha_s(\mu_b)}^{\alpha_s(\mu)} d\alpha' \frac{\Gamma_{\text{cusp}}(\alpha')}{\beta(\alpha')} \int_{\alpha_s(\mu_b)}^{\alpha'} \frac{d\alpha}{\beta(\alpha)}. \quad (20)$$

One can perform the above integration first expanding the β -function and finally re-writing $\alpha_s(\mu_b)$ in terms of $\alpha_s(\mu)$ at the correct order, as shown in Ref. [3]. The result is:

$$\begin{aligned} h_\Gamma^R(b_T; \mu) = & \frac{\Gamma_0(X - (X-1)\ln(1-X))}{2a_s\beta_0^2} \\ & + \frac{\beta_1\Gamma_0(2X + \ln^2(1-X) + 2\ln(1-X)) - 2\beta_0\Gamma_1(X + \ln(1-X))}{4\beta_0^3} \\ & + \frac{a_s}{4\beta_0^4(1-X)} (\beta_0^2\Gamma_2X^2 - \beta_0(\beta_1\Gamma_1(X(X+2) + 2\ln(1-X)) \\ & \quad + \beta_2\Gamma_0((X-2)X + 2(X-1)\ln(1-X))) \\ & \quad + \beta_1^2\Gamma_0(X + \ln(1-X))^2), \end{aligned} \quad (21)$$

while for the $\hat{C}$'s one gets:

$$\begin{aligned} \hat{C}_{q \leftarrow q'}(z, b_T; \mu) = & \left\{ \delta(1-z) + 2a_s C_F \left[1-z - \delta(1-z) \frac{\pi^2}{12} - \mathcal{P}_{q \leftarrow q'} L_T \right] \right\} \delta_{qq'} \\ \hat{C}_{q \leftarrow g}(z, b_T; \mu) = & 2a_s T_F [2z(1-z) - \mathcal{P}_{q \leftarrow g} L_T]. \end{aligned} \quad (22)$$

Thus, the fully resummed perturbative part of the TMDPDF reads:

$$\begin{aligned} \bar{F}_{q/N}^{\text{pert}}(x, b_T; \zeta, \mu) = & \left(\frac{\zeta}{\mu_b} \right)^{-D^R(b_T; \mu)} e^{h_\Gamma^R(b_T; \mu) - h_\gamma^R(b_T; \mu)} \\ & \sum_j \int_x^1 \frac{dz}{z} \hat{C}_{q \leftarrow j}(x/z, b_T; \mu) f_{j/N}(z; \mu). \end{aligned} \quad (23)$$

2. Phenomenological analysis

Eq. (23) has been obtained resumming perturbatively the large logs. Therefore, it is valid only in the perturbative region. The large $-b_T$ region of the TMDPDF must be corrected for non-perturbative contributions:

$$\tilde{F}_{q/N}(x, b_T; Q_i^2, \mu_i) = \tilde{F}_{q/N}^{\text{pert}}(x, b_T; Q_i^2, \mu_i) \tilde{F}_{q/N}^{\text{NP}}(x, b_T) \left(\frac{Q_i^2}{Q_0^2} \right)^{-D^{\text{NP}}(b_T)}, \quad (24)$$

where we set $\zeta \equiv Q^2$ and $\tilde{F}_{q/N}^{\text{NP}}(x, b_T)$ is the non-perturbative part of the TMDPDF.

We notice that the splitting into a coefficient and a PDF is valid only at high transverse momentum. In order to minimize the logarithms generated by this splitting we choose $\mu_i \equiv \sqrt{\zeta_i} \equiv Q_i = Q_0 + q_T$ where Q_0 is a fixed low scale. In our analysis we choose $Q_0 \sim 2$ GeV.

Collecting all results in Eqs. (3), (10) and (24) we can write the resummed TMDPDF that enters into the factorization theorem as

$$\begin{aligned} \tilde{F}_{q/N}(x, b_T; Q^2, Q) = & \exp \left\{ \int_{Q_i}^Q \frac{d\bar{\mu}}{\bar{\mu}} \gamma_F \left(\alpha_s(\bar{\mu}), \ln \frac{Q^2}{\bar{\mu}^2} \right) \right\} \\ & \times \left(\frac{Q^2}{\mu_b^2} \right)^{-D^R(b_T; Q_i)} e^{h_\Gamma^R(b_T; Q_i) - h_\gamma^R(b_T; Q_i)} \\ & \times \sum_j \int_x^1 \frac{dz}{z} \hat{C}_{q \leftarrow j}(x/z, b_T; Q_i) f_{j/N}(z; Q_i) \\ & \times \tilde{F}_{q/N}^{\text{NP}}(x, b_T) \left(\frac{Q_i^2}{Q_0^2} \right)^{-D^{\text{NP}}(b_T)}, \end{aligned} \quad (25)$$

where we set $\mu_f \equiv \sqrt{\zeta_F} \equiv Q$ and $\mu_i \equiv \sqrt{\zeta_i} \equiv Q_i = Q_0 + q_T$.

In our analysis we considered the low energy Drell-Yan data of the Fermilab E288⁸ experiment at three different energies of the beam (200, 300, 400 GeV/c), the CERN R209 data,⁹ and the Tevatron data of CDF Run I [10] and Run II [11] and D0 Run I [12, 13] and Run II [14]. We performed the analysis at Next-to-Leading-Logarithmic (NLL) and NNLL accuracy. Correspondingly, for the collinear PDF we used the MSTW08 PDF set¹⁵ at Next Leading Order (NLO) and Next to NLO (NNLO). Assuming $D^{\text{NP}} = 0$ in Eq. (25), we found that data can be described by the following parametrization of the non-perturbative function $\tilde{F}^{\text{NP}}$:

$$\tilde{F}_{q/N}^{\text{NP}}(x, b_T) = e^{-\lambda_1 b_T} (1 + \lambda_2 b_T^2), \quad (26)$$

where λ_1 and λ_2 are free parameters. Because CERN R209 and Fermilab E288 data are affected by large uncertainties in the normalization of the cross section we allow for two extra parameters in the fit, namely the normalization factors for E288 and R209 data: N_{E288} and N_{R209} . Therefore, we have, in total, four free parameters: λ_1 , λ_2 , N_{E288} and N_{R209} . The results of our fits are summarized in Tables (1) and (2). We found that both NLL-NLO and NNLL-NNLO calculations

Table 1. Results of our global fit on low-energy and Tevatron data, with $D^{\text{NP}} = 0$ (Eq. (26)), $Q_i = Q_0 + q_T$, at NNLL and NNL accuracies and with the collinear parton distributions from MSTW08 at NNLO and NLO.

NLL	223 points	$\chi^2/\text{d.o.f.} = 1.51$
	$\lambda_1 = 0.26^{+0.05}_{-0.02}_{\text{th}} \pm 0.05_{\text{stat}} \text{ GeV}$	$\lambda_2 = 0.13 \pm 0.01_{\text{th}} \pm 0.03_{\text{stat}} \text{ GeV}^2$
	$N_{\text{E288}} = 0.9^{+0.2}_{-0.1}_{\text{th}} \pm 0.04_{\text{stat}}$	$N_{\text{R209}} = 1.3 \pm 0.01_{\text{th}} \pm 0.2_{\text{stat}}$
NNLL	223 points	$\chi^2/\text{d.o.f.} = 1.12$
	$\lambda_1 = 0.33 \pm 0.02_{\text{th}} \pm 0.05_{\text{stat}} \text{ GeV}$	$\lambda_2 = 0.13 \pm 0.01_{\text{th}} \pm 0.03_{\text{stat}} \text{ GeV}^2$
	$N_{\text{E288}} = 0.85 \pm 0.01_{\text{th}} \pm 0.04_{\text{stat}}$	$N_{\text{R209}} = 1.5 \pm 0.01_{\text{th}} \pm 0.2_{\text{stat}}$

Table 2. Total and partial χ^2/points of our global fit on low-energy and Tevatron data with $D^{\text{NP}} = 0$ (Eq. (26)), $Q_i = Q_0 + q_T$, at NNLL and NNL accuracies and with the collinear parton distributions from MSTW08 at NNLO and NLO.

		NNLL, NNLO	NNL, NLO
	points	χ^2/points	χ^2/points
	223	1.10	1.48
E288 200	35	1.53	2.60
E288 300	35	1.50	1.12
E288 400	49	2.07	1.79
R209	6	0.16	0.25
CDF Run I	32	0.74	1.31
D0 Run I	16	0.43	1.44
CDF Run II	41	0.30	0.62
D0 Run II	9	0.61	2.40

give a good description of the data. In particular NNLL-NNLO calculation gives a better total χ^2 , with a better description of Tevatron data. Notice that high-energy Tevatron data correspond to a region where $Q \equiv M_Z$ and q_T spans a region from one up to tens of GeVs. This means that, in the Fourier impact parameter space, we are exploring the small b_T region, that is, the region where the (perturbative) resummation of large logarithms plays an important role. Comparing the first and the second column of Tab. (2) we can see that, as expected, the description of Tevatron data improves going from NLL to NNLL accuracy. In Ref. [4] we explored the role played by the collinear PDFs adopting another set, namely the CTEQ10 [16] and we find a complete consistency among the results. We also checked the role played by a possible D^{NP} function different from zero. We did not find a conclusive evidence of such a term in present DY/Z data. However this term may be relevant for very low energy data sets like the SIDIS data sets of the HERMES and COMPASS experiments.

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