Comparing Resummations

Workshop on Resummation and Parton Showers
IPPP Durham, July 16, 2013

Based on ongoing work with Mao Zeng (to appear) with a nod to
E. Laenen, GS, W. Vogelsang, Phys. Rev. D63 114018 (2001)
and even earlier work.

A. Threshold resummation

B. Unity and difference: "D(irect)QCD" & "S(C)ET" evolution-based threshold resummations

C. Diagrammatic resummation for electroweak annihilation and power corrections

D. DQCD and SCET for hadrons

Some relatively simple observations that hopefully won't get lost in the formulas.

A. Threshold Resummation

- **Electroweak boson production** $(E = Z, W, H)$ **mass** M**;**

$$A + B \;\to\; E(M) + X$$
$$\tau \;\equiv\; \frac{M^2}{S} \quad \text{true threshold : } S \to M^2 \;\leftrightarrow\; \tau \to 1$$

- **Collinear factorization:**

$$\frac{d\sigma_{AB\to M}(S, M^2)}{dM^2} = \sum_{\text{partons }a,b} \frac{\Sigma_{ab}}{S\,Q^2} \int_0^1 \frac{dx_a}{x_a}\frac{dx_b}{x_b}\; f_{a/A}(x_a,\mu_f)\, f_{b/B}(x_b,\mu_f) \times C_{ab\to E}\left(M^2/\hat{s}, M^2/\mu_f^2, \alpha_s(\mu_f)\right)$$

$$= \sum_{\text{partons }a,b} \frac{\Sigma_{ab}}{S\,Q^2} \int_\tau^1 \frac{dz}{z}\; \mathcal{L}_{ab}\left(\frac{\tau}{z},\mu_f\right)\; C_{ab\to E}\left(z, M^2/\mu_f^2, \alpha_s(\mu_f)\right)$$

- **Partonic luminosity:**

$$\mathcal{L}_{ab}\left(\frac{\tau}{z},\mu_f\right) = \int_0^1 \frac{dx_a}{x_a}\frac{dx_b}{x_b}\; f_{a/A}(x_a,\mu_f)\, f_{b/B}(x_b,\mu_f)\,\delta\left(1 - \frac{\tau/z}{x_a x_b}\right).$$

in terms of

$$z \equiv \frac{\tau}{x_a x_b} \quad \text{partonic threshold : } x_a x_b S \to M^2 \;\leftrightarrow\; z \to 1$$

- **Partonic hard-scattering functions are generally singular at partonic threshold for $b = \bar{a}$: (purely gluonic radiation)**

$$C_{a\bar{a}}^{(n)} \sim \alpha_s{}^n(\mu_f) \left[\frac{\ln^{2n-1}(1-z)}{1-z}\right]_+ , \; \alpha_s{}^n(\mu_f) \left[\frac{\ln^{2n-2}(1-z)}{1-z}\right]_+ ,$$

$$\dots \quad \alpha_s{}^n(\mu_f)\, \delta(1-z)$$

- **In hadronic scattering these are enhancements, due to incomplete cancellation between virtual and radiated gluons in the cross section *and* in subtraction of soft/collinear gluons.**

- **(Nearly) two logs per loop: collinear and soft divergences**

- **Threshold resummation organizes singular distributions for $z \to 1$. Clearly a good idea.**

- **The solution turns out to be simple, and based on factorization:**

$$C_{ab\to M}\left(z, M/\mu_f, \alpha_s(\mu_f)\right) = \delta_{b\bar{a}} \; H_{a\bar{a}}\left(M, \mu_f, \alpha_s(\mu_f)\right) \times S_{a\bar{a}}(M(1-z)/\mu_f, \alpha_s(\mu_f))$$

$$+\mathcal{O}\left(1-z\right)$$

- **$S(M(1-z))$ is the "soft function": common to DQCD and SCET analyses**

- Kinematics shows why they're the so similar. For $z \to 1$, let k_0 be total c.m. energy of radiation, then

$$1 - z = \frac{2k_0}{M}$$

- For $z \to 1$ all radiation is in the soft sector, the "S" of SCET.
The collinear Lagrangian doesn't play as important role in the z-dependence.

- The Soft function is a DQCD collinear-subtracted $\leftrightarrow$ SCET renormalized matrix element.
(GS Nucl.Phys. B281 (1987) G. Marchesini, G. Korchemsky Nucl.Phys. B406 (1993) 225, A. Belitsky Phys.Lett. B442 (1998) 307, T. Becher, M. Neubert, G. Xu JHEP 0807 (2008) 030)

$$S_{a\bar{a}}(2k_0/\mu_f) = M \int \frac{dx^0}{4\pi} e^{ik_0 x^0} W_{\rm DY}(x^0 \mu_f, \vec{x} = 0)$$

$$W_{a\bar{a}}(x\mu_f) = \frac{1}{N_c} \langle 0| \,\mathrm{Tr}\, \bar{\mathrm{T}}[\Phi_n^\dagger(x)\Phi_{\bar{n}}(x)]\, \mathrm{T}[\Phi_{\bar{n}}^\dagger(0)\Phi_n(0)]|0\rangle$$

in terms of annihilating, lightlike Wilson lines (in a, $\bar{a}$ color reps):

$$\Phi_n(x) = \mathrm{P}\, \exp\left(ig \int_{-\infty}^0 ds\, n \cdot A(x + sn)\right)$$

- For direct photons, similar considerations apply
(T. Becher & M. Schwartz JHEP 1002 (2010) 040)

- At $z = 1$ infrared divergences are not cancelled.

- Singular $z \to 1$ behavior is organized by transforms

- Mellin, or Laplace

$$\widetilde{S}_{a\bar{a}}\left(\ln \frac{M^2}{\tilde{N}^2\mu^2}, \alpha_s(\mu)\right) = \int_0^1 dz\, z^{N-1} S_{a\bar{a}}\left(M(1-z), \mu\right)$$
$$= \int_0^\infty d\omega\, \exp\left(-\frac{N\omega}{M}\right) S_{a\bar{a}}\left(\omega, \mu\right) + \mathcal{O}(1/N)$$

- Correspondence:

$$\left[\frac{\ln^k(1-z)}{1-z}\right]_+ \leftrightarrow \ln^{k+1} N$$

- The problem is posed in the same way in DQCD, SCET.

- They both start with a common, evolution-based resummation, whose basis is in the factorization of soft gluon radiation.

B. Unity and Difference: DQCD & SCET Evolutions

- The N dependence of $\tilde{S}_{a\bar{a}}\left(\ln \frac{M^2}{\bar{N}^2\mu^2}, \alpha_s(\mu_f)\right)$ is determined by its μ_f dependence:

$$\frac{d}{d\ln\mu} \ln \widetilde{S}_{a\bar{a}}\left(\ln \frac{M^2}{\bar{N}^2\mu^2}, \mu\right) = -2\Gamma_{\rm cusp}\left(\alpha\right) \ln \frac{M^2}{\mu^2 \bar{N}^2} - \Gamma_{\rm DY}\left(\alpha\right)$$

- In terms of the "cusp anomalous dimension", $\Gamma_{\rm cusp}$,and $\Gamma_{\rm DY} \leftrightarrow \gamma_W$

  (A common notation $\bar{N} = Ne^{\gamma_E}$ simplifies notation at NLL)

- An equivalent momentum space evolution equation for the log is
  (GS, 1987, Belitsky 2000, Becher, Neubert 2006)

$$\frac{d}{d\ln\mu} \ln \tilde{S}_{a\bar{a}}\left(M(1-z)/\mu, \alpha(\mu)\right) = \left[-4\Gamma_{\rm cusp}\left(\alpha(\mu)\right) \ln \frac{M}{\mu} - 2\Gamma_{\rm DY}\left(\alpha(\mu)\right)\right] \delta(1-z) \\ -4\Gamma_{\rm cusp}\left(\alpha(\mu)\right) \left(\frac{1}{1-z}\right)_+$$

  The (simpler) moment-space solution:

$$\tilde{S}\left(\frac{M}{\tilde{N}\mu_1}, \alpha_s(\mu_1)\right) = \tilde{S}\left(\frac{M}{\tilde{N}\mu_2}, \alpha_s(\mu_2)\right) \exp\left\{\int_{\mu_2}^{\mu_1} \frac{d\rho}{\rho} \left(2\Gamma_{\rm cusp}(\alpha_s(\mu) \ln\left(\frac{\rho N}{M}\right) + \Gamma_{\rm DY}\left(\alpha_s(\rho)\right)\right)\right\}$$

- It is at this stage that SCET and QCD methods part ways, in their choices of the lower scale, μ_1, and then in their use of the inverse Laplace or Mellin transform to derive a physical cross section.

- **DQCD: run up from $\mu_2 = M/\bar{N}$ to $\mu_2 = \mu_f \sim M$, then absorb remaining N-dependence from the lower scale into the exponent**

$$\tilde{S}\left(\frac{M}{N\mu}, \alpha_s(\mu)\right) = S\left(1, \alpha_s(M/\tilde{N})\right) \exp\left\{\int_{M/\tilde{N}}^{\mu} \frac{d\rho}{\rho} \left(2\Gamma_{\rm cusp} \ln\left(\frac{\rho\tilde{N}}{M}\right) + \Gamma_{\rm DY}\left(\alpha_s(\rho)\right)\right)\right\}$$

$$= \tilde{S}\left(1, \alpha_s(\mu)\right) \exp\left\{\int_{M/\tilde{N}}^{\mu} \frac{d\rho}{\rho} \left(2\Gamma_{\rm cusp}(\alpha_s(\rho)) \ln\left(\frac{\rho\tilde{N}}{M}\right) + \hat{D}\left(\alpha_s(\rho)\right) \right)\right\}$$

- **Here, define $\hat{D}$ by**

$$\hat{D}\left(\alpha_s(\rho)\right) = \Gamma_{\rm DY}\left(\alpha_s(\rho)\right) - \rho\frac{\partial}{\partial\rho} \ln S(1, \alpha_s(\rho))$$

- **The function $\hat{D}$ is thus a hybrid function, the sum of an anomalous dimension and a perturbative coefficient.**

- **A conventional expression is**

$$\tilde{S}\left(\frac{M}{N\mu},\alpha_s(\mu)\right) = \tilde{S}\left(1,\alpha_s(M)\right)\ \exp\left\{\int_{1/\bar{N}}^{1}\frac{dy}{y}\left[\,2\int_{yM}^{M}\frac{dq}{q}\,\Gamma_{\rm cusp}(\alpha_s(q)) + \hat{D}\left(\alpha_s(yM)\right)\,\right]\right\}$$

- **There is a (Landau) singularity for large $N \sim M/\Lambda_{\rm QCD}$. This can be avoided at the cost of making a choice at the level of power corrections in $\Lambda_{\rm QCD}/M$.**

- **The "minimal" (CMNT) prescription (Catani, Mangano, Nason, Tretadue Nucl.Phys. B478 (1996) 273-310) can be used to avoid the singularity in computing hadronic cross sections and designed specifically to avoid introducing nonperturbative parameters or explicit power corrections. Other prescriptions are available: 'Borel' (Forte and Ridolfi) and 'principle value' (Contopanagos, GS, Contopanagos, Berger)**

- For the SCET, evolve from μ_f down to a fixed "soft" scale, μ_s

$$\tilde{S}\left(\frac{M}{\tilde{N}\mu_f}, \alpha_s(\mu_f)\right) = \tilde{S}\left(\frac{M}{\bar{N}\mu_s}, \alpha_s(\mu_s)\right) \exp\left\{\int_{\mu_f}^{\mu_s} \frac{d\rho}{\rho} \left(2\Gamma_{\rm cusp} \ln\left(\frac{\rho\bar{N}}{M}\right) + \Gamma_{\rm DY}\left(\alpha_s(\rho)\right)\right)\right\}.$$

- N-dependence is in the prefactor, the soft function evaluated at scale μ_s and the explicit factor of $\ln N$ that multiplies the cusp anomalous dimension in the exponent.

- Reorganize by separating all the $\ln N$ dependence in the exponent as a multiplicative factor, using it as a generating functional for N-dependence in the soft function at μ_s.

- Gives:

$$\begin{aligned}\widetilde{S}_{a\bar{a}}\left(\ln\frac{M^2}{\bar{N}^2\mu^2}, \mu\right) &= \exp\left[-4\mathcal{S}\left(\mu_s,\mu\right) + 2\alpha_{\gamma_{\rm DY}}\left(\mu_s,\mu\right) + \eta\,\ln\frac{M^2}{\mu_s^2}\right] \\ &\quad\times \widetilde{S}_{a\bar{a}}\left(\ln\frac{M^2}{\mu_s^2} + \frac{\partial}{\partial\eta}, \mu_s\right)\,\exp\left[-\,\eta\,\ln\bar{N}^2\right]\end{aligned}$$

with

$$\begin{aligned}\mathcal{S}\left(\mu_s,\mu\right) &= -\int_{\mu_s}^{\mu} d\ln\mu'\,\Gamma_{\rm cusp}(\alpha_s(\mu))\,\ln\frac{\mu'}{\mu_s} \\ \alpha_{\gamma_{\rm DY}}\left(\mu_s,\mu\right) &= -\int_{\mu_s}^{\mu} d\ln\mu'\,\Gamma_{\rm DY}(\alpha_s(\mu))\end{aligned}$$

and

$$\eta = 2\int_{\mu_s}^{\mu} d\ln\mu'\,\Gamma_{\rm cusp}(\alpha)$$

- In EFT, the soft scale μ_s is to be chosen to stabilize the cross section when $S_{a\bar{a}}$ is known to low order. The assumption is that the range of z for which $(1-z)M < \mu_s$ is not phenomenologically important, even though it is full of logs of $1-z$. This is termed "dynamical scale generation"

- Not too surprisingly (given their origin) the two formalisms agree to non-leading logarithms in N at fixed N (Bonvini, Forte, Marzani, Ridolfi 1303.3590 & Nucl.Phys. B861 (2012) 337-360).

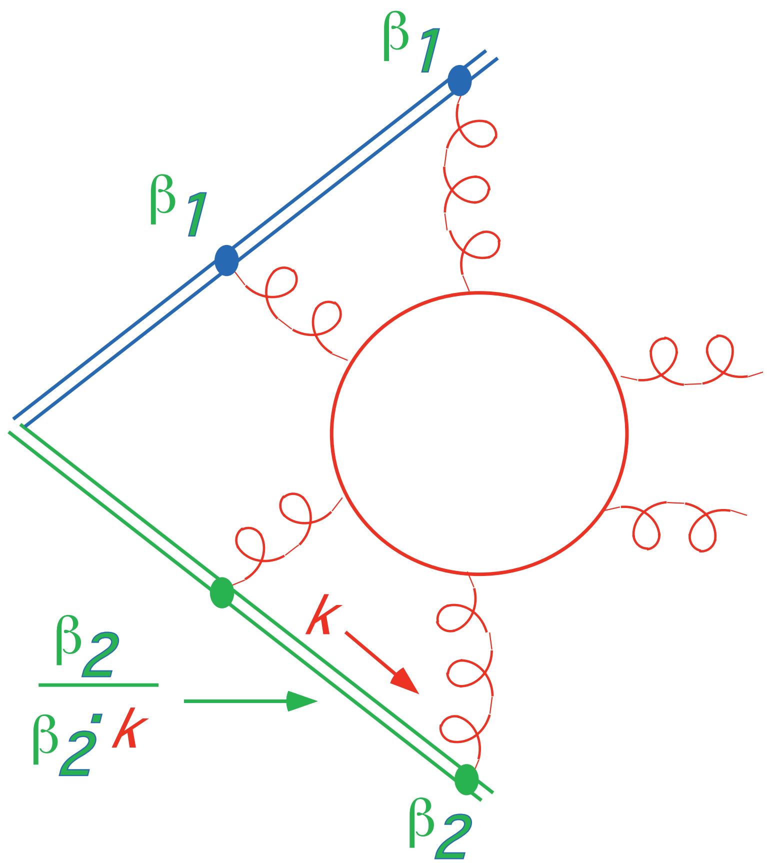
- The EFT resummation (Becher, Neubert, Schwartz, Stewart . . . 2006 - . . .) is even more insensitive to the Landau pole than the minimal prescription in its application to hadronic cross sections, but at the price of abandoning non leading logarithmic dependence

- In a sense, the DQCD minimal prescription tries to "have it both ways", all orders in N in moment space, and no new nonperturbative parameters. But this is a particular choice of nonperturbative parameters.

- We'll return to their comparison below, but first have a look at the nonperturbative limit both formalisms are designed to avoid.

C. Diagrammatic resummation and power corrections

- For products of Wilson lines, exponentiation is exact in terms of “webs” (Gardi, Smille, White, Laenen 2009 ... 2013)
- Look for further insight into the “independent emission” – exponentiation – by analogy to soft photons.
- A typical diagram (for final-state sources)



- **Webs and exponentiation for soft contributions to weighted ($e = Q_T$, thrust, $and\ k_0$) cross sections** (GS (1981); Gatheral & Frenkel and Taylor, (1981), Mitov, GS, Sung (2009)). **A convolution,**

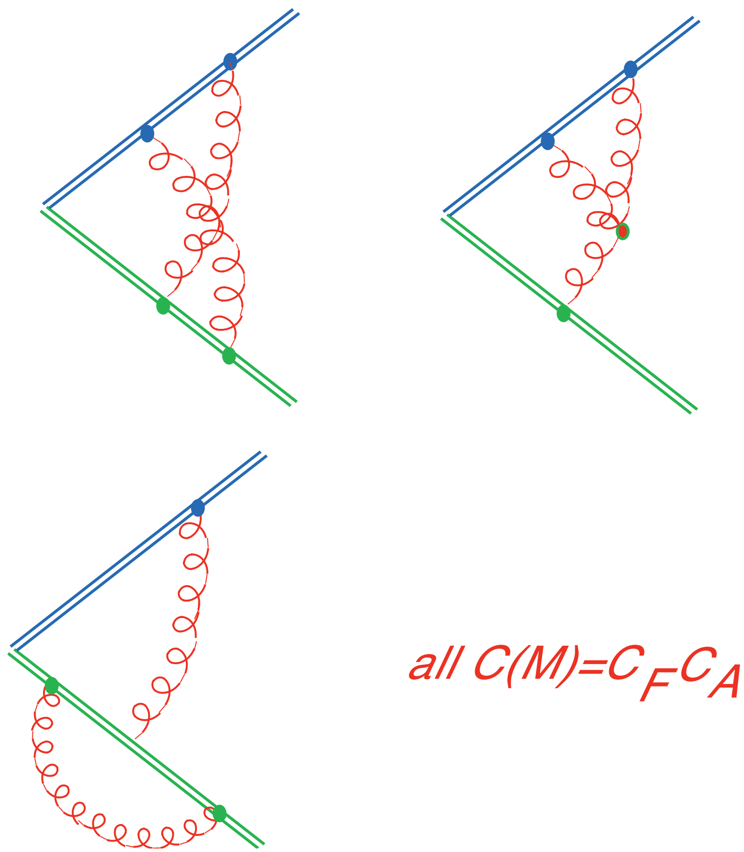
$$\frac{d\sigma}{de_a} = \sum_{n=0}^{\infty} \frac{1}{n!} \int de\, \delta(e - e_s)\ \otimes_{\sum e_i = e_a} \prod_{i=1}^{n} E(e_i)$$

- **With a kernel that becomes the exponent in transform space**

$$E(e) = \sum_{\text{states n}} \mathcal{W}_n(e) = \sum_{\mathcal{M}} C(\mathcal{M}_n)\, \mathcal{M}_n^2(e)$$

- **The $\mathcal{M}_n^2$ are the corresponding momentum integrals for the web diagrams.**

- **The $C(\mathcal{M}_n)$ are modified color factors for $\mathcal{M}_n$s. Examples at α_s^2:**



- **Notice that non-planar diagrams contribute in $N_c \to \infty$ limit**

- The webs determine directly determine the exponent under transforms:

$$\tilde{S}_e(N) \equiv \int de\ \mathrm{e}^{-Ne}\ \frac{dS}{de} = \exp\left[\int de'\ \mathrm{e}^{-Ne'}\ E(e')\right]$$

- Double logarithmic behavior is encoded in the construction of the webs $\mathcal{W}$. Subdivergences cancel.

- Each web gives a single collinear and infrared logarithm up to running coupling effects, just like a single gluon.

- In a theory with a fixed coupling (SYM . . .) a web acts exactly like a single gluon.

- The 2-loop structure of Γ_S is an intriguing suggestion that “web=NP gluon” could generalize to arbitrary hard processes.

- For electroweak annihilation as above, this gives a very specific template for the all-orders form, up to exponentially-suppressed corrections in N:

- Boost invariance along dipole axis $\Rightarrow$
(Laenen, GS, Vogelsang, 2001)

$$\ln C(N,M) = \sum_N \int d\mathrm{PS}_N \, \theta(M^2 - k^2)|M^{(\mathrm{eik})}|^2 \, e^{-N(2k_0/M)}$$

$$= \int_0^{M^2} \frac{\rho(\alpha_s(u,\varepsilon))}{u^2} \left[K_0\left(\frac{2Nu}{M}\right) + \ln\frac{u}{M}\right] + \ln\bar{N} \int_0^{M^2} \frac{du^2}{u^2} A(\alpha_s(u,\varepsilon))$$

- The "α_s" for a single gluon is replaced by:

$$\rho(\alpha_s(u,\varepsilon)) = \Gamma_{\mathrm{cusp}}(\alpha_s(u,\varepsilon)) + \frac{\partial D'}{\partial \ln \mu^2}$$

- $\Gamma_{\mathrm{cusp}} =$ "A" is the same "cusp" anomalous dimension, and $D'(\alpha_s)$ is analogous to the D-term as above.

- The structure of "renormalon" power corrections is given by the low-x expansion of $K_0(x) - \ln(x)$

- Only even powers of $\frac{N}{M}$ occur (as noted for "large-N_f" by Beneke and Braun (1995))

- This phenomenology has yet to be developed for hadron-hadron scattering, is (mercifully) irrelevant to inclusive Higgs etc. at the LHC, but for large-N_c, all radiation is from dipoles.

D. Cross sections for hadrons

- **Does a quantifiable correspondence in moment space breaks down for hadronic cross sections.?** (Bonvini, Forte, Ghezzi, Ridolfi, 2012,13) **I'll argue here it's "natural" for the two evolution-based resummations to return very similar phenomenological results, at least when the corrections are significant.**
- **The differing expressions for the cross section**
- **DQCD as an inverse transform, to be performed numerically in general**

$$\frac{d\sigma_{AB\to M}(S,M^2)}{dM^2} = \sum_{\text{partons } a,b} \frac{\Sigma_{ab}}{S\,Q^2} \int_0^1 \frac{dx_a}{x_a}\frac{dx_b}{x_b}\, f_{a/A}(x_1,\mu_f)\, f_{b/B}(x_2,\mu_f) \times C_{ab\to E}\left(M^2/\hat{s}, M^2/\mu_f^2, \alpha_s(\mu_f)\right)$$

$$= \sum_{\text{partons } a,b} \frac{\Sigma_{ab}}{S\,Q^2} \int_\tau^1 \frac{dz}{z}\, \mathcal{L}_{ab}\left(\frac{\tau}{z},\mu_f\right)\, C_{ab\to E}\left(z, M^2/\mu_f^2, \alpha_s(\mu_f)\right)$$

- **which can be written as**

$$\frac{d\sigma_{AB\to E}(\tau,M^2)}{dM^2} = \sum_a H_{a\bar{a}}(M,\mu_f) \int_{n_0-i\infty}^{n_0+i\infty} \frac{dN}{2\pi i}\, \tau^{-N}\, \mathcal{L}_{a\bar{a}}(N,\mu_f) \times \exp\left\{ \int_{1/\bar{N}}^1 \frac{dy}{y}\left[2\int_{yQ}^{Q} \frac{dq}{q}\, \Gamma_{\text{cusp}}(\alpha_s(q)) + \hat{D}\left(\alpha_s(yQ)\right) \right]\right\}$$

- **SCET: The inverse transform of the explicit N-dependence can be done exactly,**

$$
\begin{aligned}
S_{a\bar{a}}\left(\ln\frac{M^2(1-z)^2}{\mu^2},\alpha_s(\mu)\right) &= \exp\left[-4S\left(\mu_s,\mu\right)+2\alpha_{\gamma_{\rm DY}}\left(\mu_s,\mu\right)+\eta\,\ln\frac{M^2}{\mu_s^2}\right] \\
&\quad\times\,\widetilde{S}_{\rm DY}\left(\ln\frac{M^2}{\mu_s^2}+\frac{\partial}{\partial\eta},\mu_s\right)\;(1-z)^{2\eta-1}\frac{1}{\Gamma(2\eta)}
\end{aligned}
$$

- **So that the cross section is given directly as**

$$
\begin{aligned}
\frac{d\sigma_{AB\to M}(S,M^2)}{dM^2} &= \int_\tau^1\frac{dz}{z}\,\mathcal{L}_{ab}\left(\frac{\tau}{z},\mu_f\right) \\
&\quad\times H_{a\bar{a}}(M/\mu,\alpha_s(\mu_f))\;S_{a\bar{a}}\left(\ln\frac{M^2(1-z)^2}{\mu^2},\alpha_s(\mu)\right)
\end{aligned}
$$

- **Which is also adaptable to differential cross sections.**

- **The 'problem' in short:**

- **In DQCD all N-dependence is kept, even though the inverse transform is then influenced somewhat by a branch cut at large, real N.**

- **In SCET, higher-order N dependence is sacrificed to get a well-defined transform.**

- **Why/when should they return the comparable answers?**

- If the parton luminosity were exactly power-behaved ("single-power approximation")

$$\mathcal{L}_{a\bar{a}}\left(\frac{\tau}{z}\right) = \text{const}\ \left[\frac{\tau}{z}\right]^{-s(\tau)}$$

$$= \mathcal{L}(\tau)\, z^{s(\tau)}$$

the convolution would itself be a moment of the partonic cross section:

$$\frac{d\sigma(\tau)}{dM^2} = \sigma_0 \int_\tau^1 \frac{dz}{z}\mathcal{L}\left(\frac{\tau}{z}\right)\omega(z)$$

$$= \sigma_0\ \mathcal{L}(\tau) \int_\tau^1 \frac{dz}{z} z^{s(\tau)}\omega(z)$$

$$= \sigma_0\ \mathcal{L}(\tau)\ \widetilde{\omega}\left((s(\tau)\right)$$

- If this were the case, then when variant resummations for ω agree at the level of moments, they should agree phenomenologically.
- And, perhaps surprisingly, this happens in a wide range of realistic cases.

- **The observation: if $\mathcal{L}(\tau/z)$ falls rapidly as a function of z, the single-power approximation will hold for the cross section if $\mathcal{L}(\tau/z)$ becomes negligible before $n \geq 2$ terms in the expansion**

$$\ln \mathcal{L}\left(\frac{\tau}{z}\right) = \sum_{n=0} \frac{1}{n!} s_n(\tau) \left(\ln \frac{1}{z}\right)^n$$

become important, with $s_n = d^n \ln \mathcal{L}(\tau')/d \ln^n \tau'|_{\tau'=\tau}$.

- **Since the decay distance in $ln\frac{1}{z}$ is $s_1(\tau)$, we need**

$$\frac{1}{n!} \frac{s_n(\tau)}{[s_1(\tau)]^n} \ll 1$$

- **For example, model parton distributions like**

$$f(x) = C\, x^{-\gamma}\, (1-x)^{\beta}$$

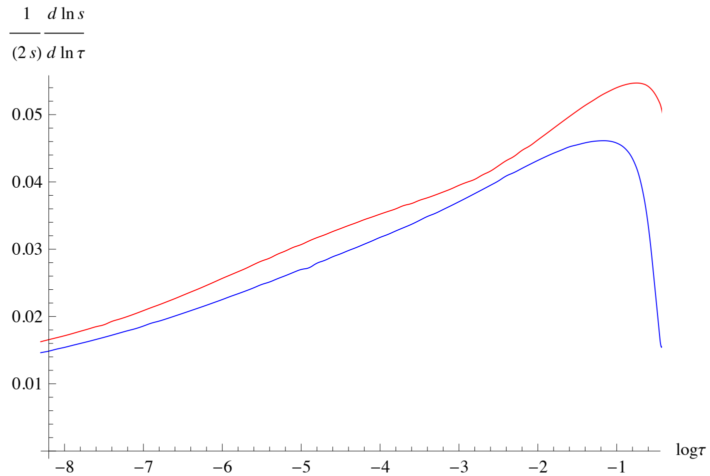
give luminosity functions (Appel, Mackenzie, GS Nucl.Phys. B309 (1988) 259, Becher, Neubert JHEP 0807 (2008) 030)

$$\mathcal{L}(\tau') = C^2\, B(\beta+1, \beta+1)\, \tau'^{-\gamma}\, (1-\tau')^{2\beta+1}\, F_{2,1}\left(\beta+1, \beta+1; 2\beta+2, 1-\tau'\right)$$

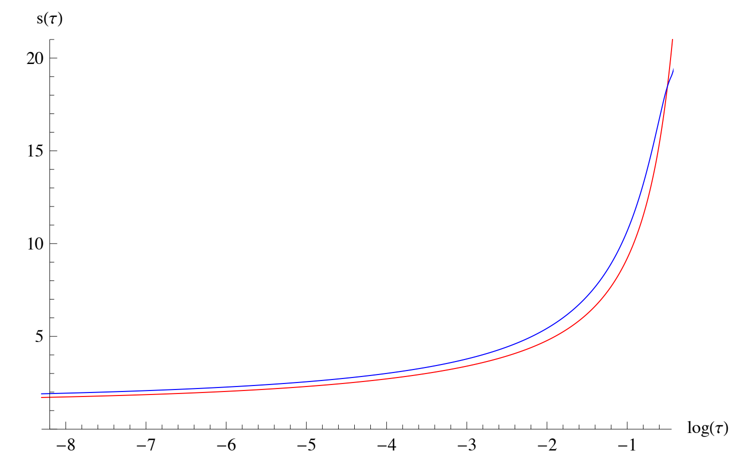
- **For this explicit form one can show: $\mathcal{L}$ is single-power dominated:**

1. **for small τ even for moderate $\gamma = \mathcal{O}(1)$**
2. **for all τ for $\beta \to \infty$**

- **And for “real” PDFs (MRST2008 NNLO gluon) at M=10 GeV (red) and M=126 GeV (blue)**



- **To be compared to**



- **the first derivative dominates the second over an adequate region, even when the first derivative is not very large.**

- So, for many, probably most cases, the effective power $s_1(\tau)$ determines the dominant moment, and consequently an effective "soft scale", $\mu_s = M/s_1(\tau)$.

- And we should expect both SCET and DQCD results to be similar, with differences due primarily to "incidental" choices of scales.

Summary

- Commonly-used effective theory and full QCD threshold resummations are based on the same underlying factorization, and differ mainly in how the resulting evolution equations are solved.

- Although both formalisms have been designed to avoid power corrections, such corrections to hard-scattering functions are theoretically accessible in a manner similar to power corrections in event shapes.

- Despite their differences, DQCD and SCET resummations can be expected to give similar phenomenological predictions for a wide range of kinematics and parton distributions.