Summations of large logarithms by parton showers

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Perturbative cross sections

The main focus of this workshop is to calculate the pQCD cross sections as precise as possible, thus we have a pretty integral

$$\sigma[O_J] = \sum_m \frac{1}{m!} \sum_{\{a,b,f_1,\dots,f_m\}} \int_0^1 d\eta_a \overbrace{\int_{\eta_a}^1 \frac{dz}{z}\, \Gamma^{-1}_{aa'}(z,\mu^2)\, f_{a'/A}(\eta_a/z,\mu^2)}^{\text{Bare PDF}}$$

$$\times \int_0^1 d\eta_b \int_{\eta_b}^1 \frac{d\bar{z}}{\bar{z}}\, \Gamma^{-1}_{bb'}(\bar{z},\mu^2)\, f_{b'/A}(\eta_b/\bar{z},\mu^2)$$

$$\times \int d\phi(\eta_a\eta_b s, \{p,f\}_m) \left\langle M(\{p,f\}_m) \right| \underbrace{O_J(\{p,f\}_m)}_{\text{IR safe measurement operator}} \overbrace{\left| M(\{p,f\}_m) \right\rangle}^{\text{Partonic matrix element}}$$

$$+ \mathcal{O}\left(\frac{\Lambda^2_{QCD}}{\mu_J^2}\right)$$

Error of the factorization
(Cannot be beaten by calculating higher and higher order.)

and here the MSbar parton in parton renormalised PDF is

$$\Gamma_{aa'}(z,\mu^2) = \delta(1-z)\delta_{aa'} - \frac{\alpha_s(\mu^2)}{2\pi} \frac{1}{\epsilon} \frac{(4\pi)^\epsilon}{\Gamma(1-\epsilon)} P_{aa'}(z) + \cdots \quad .$$

Statistical space

Introducing the statistical space we can represent the QCD density operator as a vector

Bare PDFs for both incoming hadrons

$$\sigma[O_J] = \underbrace{(1|}_{} \; \mathcal{O}_J \; \overbrace{\left[\mathcal{F}(\mu^2) \circ \mathcal{Z}_F(\mu^2)\right]}^{} \; \underbrace{|\rho(\mu^2))}_{|M\rangle\langle M|}$$

All the initial and final state sums and integrals

QCD density operator
Describes the fully exclusive partonic final states.

The physical cross section is RG invariant as well as the QCD density operator and the bare PDF.

$$\mu^2 \frac{d}{d\mu^2} |\rho(\mu^2)) = \mu^2 \frac{d}{d\mu^2} \left[\mathcal{F}(\mu^2) \circ \mathcal{Z}_F(\mu^2)\right] = 0 + \mathcal{O}(\alpha_s^{k+1})$$

Perturbative expansion of the density operator

$$|\rho(\mu^2)) = \sum_{n=0}^{k} \left[\frac{\alpha_s(\mu^2)}{2\pi}\right]^n \sum_{\substack{n_R=0 \\ n_R+n_V=n}}^{n} \sum_{n_V=0}^{n} |\rho^{(n_R,n_V)}(\mu^2))$$

Number of real radiations

Number of loops

Statistical space

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A vector in the statistical space can be translated as

$$\big(\{p,f,c,s,c',s'\}_m\big|\rho\big) \Longleftrightarrow \big\langle\{c,s\}_m\big|M(\{p,f\}_m)\big\rangle\big\langle M(\{p,f\}_m)\big|\{c',s'\}_m\big\rangle$$

An operator in the statistical space corresponds to a direct products of the corresponding quantum operators:

$$\mathcal{A}(\mu^2) \Longleftrightarrow A^L(\mu^2)\otimes A^R(\mu^2)^\dagger$$

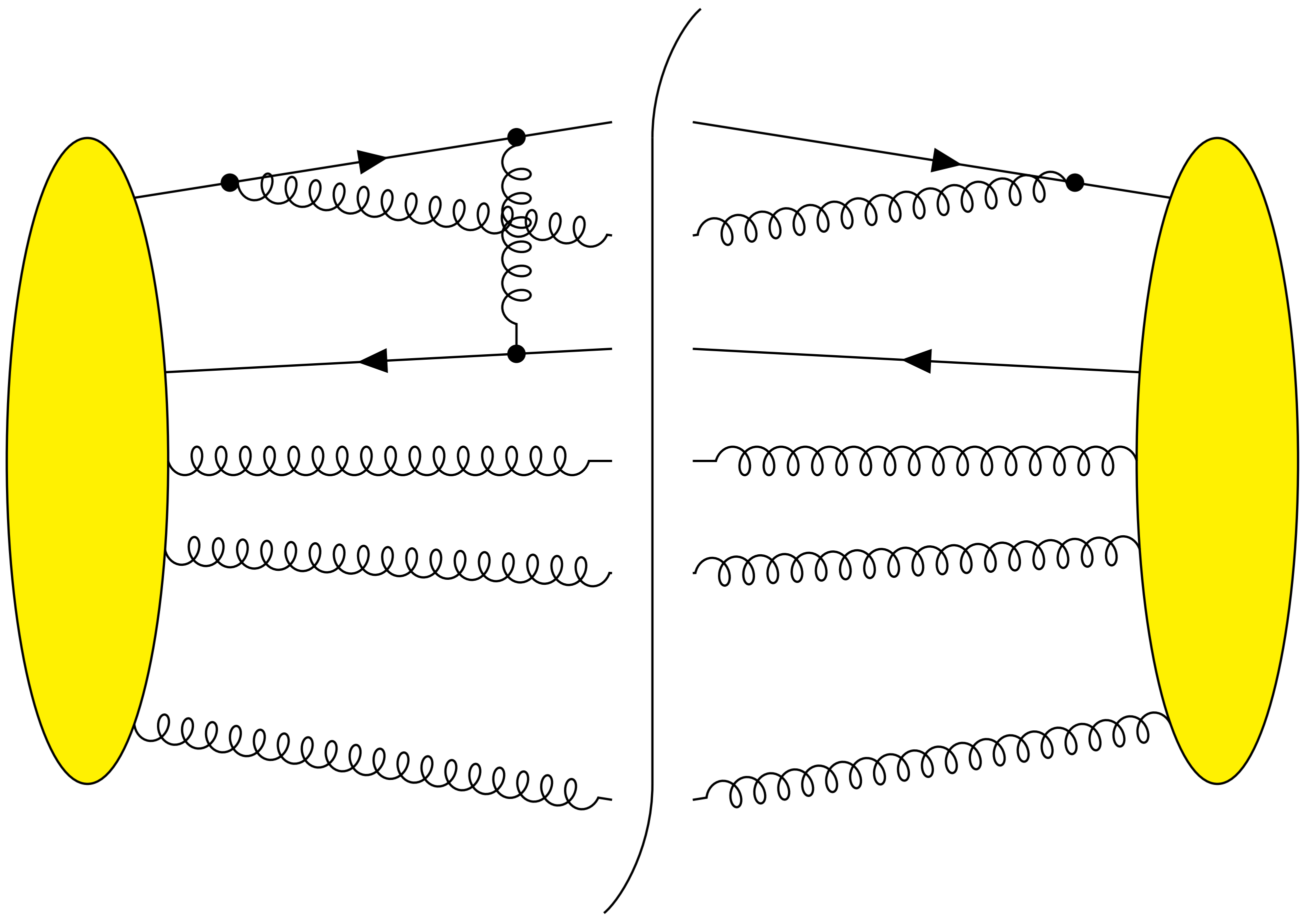
When operators act on a state we have

$$\cdots\mathcal{A}_3(\mu_3^2)\mathcal{A}_2(\mu_2^2)\mathcal{A}_1(\mu_1^2)\big|\rho\big) \Longleftrightarrow \cdots A_3^L(\mu_3^2)A_2^L(\mu_2^2)A_1^L(\mu_1^2)\big|M\big\rangle\big\langle M\big|A_1^R(\mu_1^2)^\dagger A_2^R(\mu_2^2)^\dagger A_3^R(\mu_3^2)^\dagger\cdots$$

Fixed order cross sections

Infrared sensitive operator

Amplitudes have soft or collinear singularities and they have divergences $1/\varepsilon$ from the loops

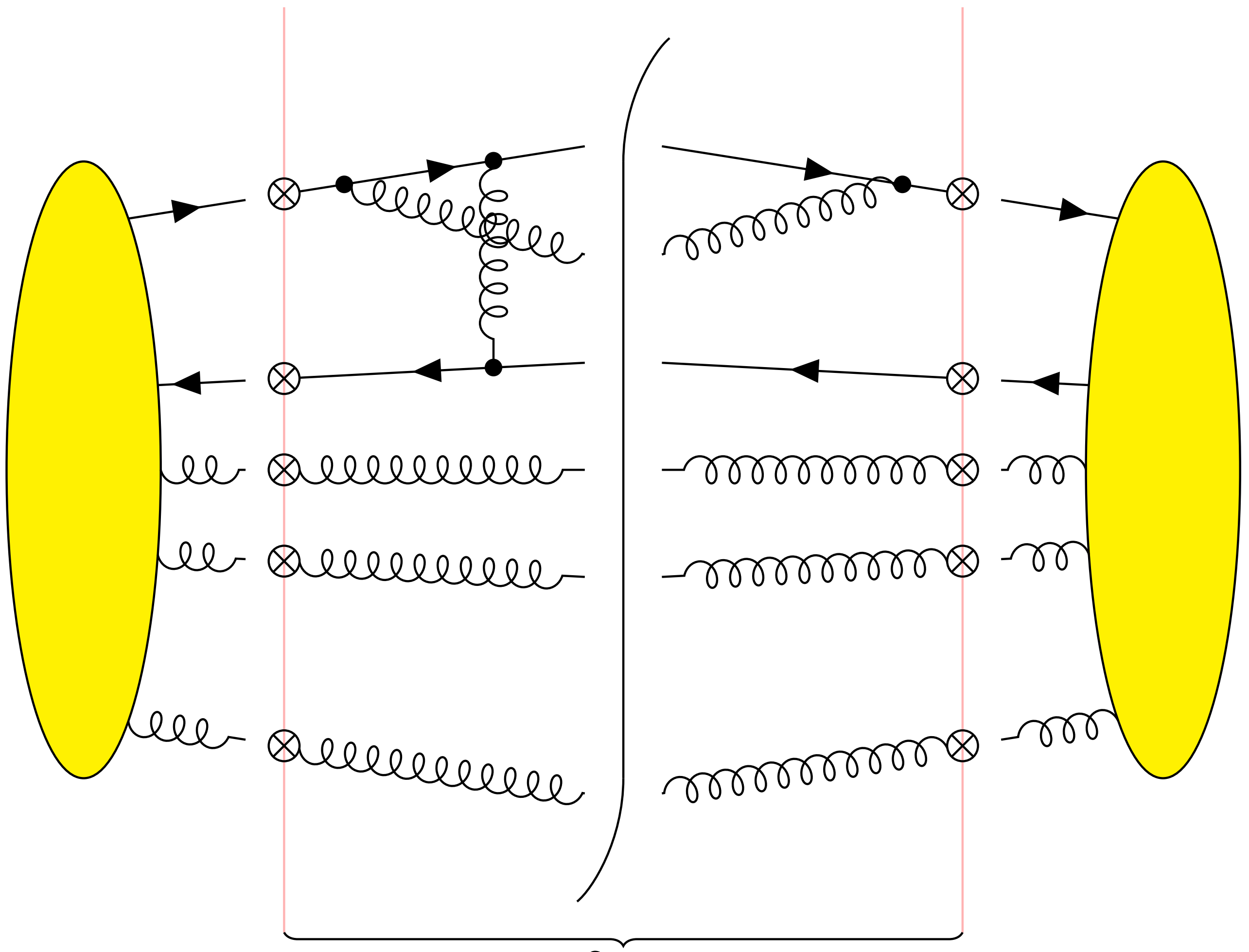


➠ We want to describe the singularity structure in **process independent way**.

➠ Everything in the yellow blobs is considered hard.

Infrared sensitive operator

Consider the momenta coming from the hard part as fixed and on shell.



$$\left[\frac{\alpha_s(\mu^2)}{2\pi}\right]^2 \mathcal{D}^{(1,1)}(\mu^2)$$

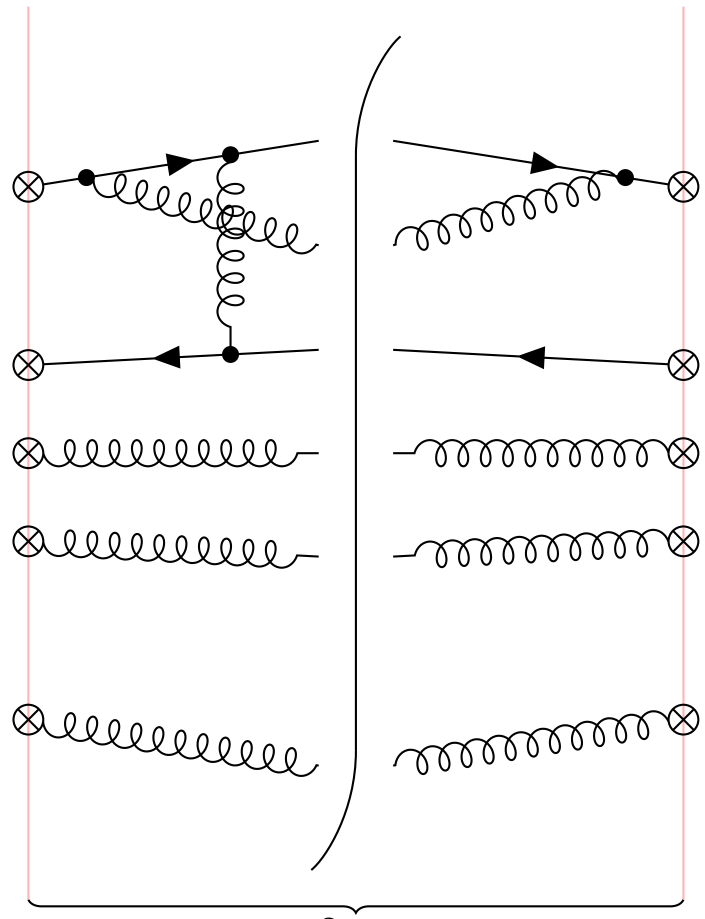
This gives us an operator as

$$
\begin{aligned}
&\big(\{\hat p, \hat f, \hat s, \hat s', \hat c, \hat c'\}_{m+n_{\rm R}}\big|\rho(\mu^2)\big) \\
&\quad \sim \frac{1}{m!}\int[d\{p\}_m]\sum_{\{f\}_m}\sum_{\{s,s',c,c'\}_m} \\
&\qquad \times \big(\{\hat p, \hat f, \hat s, \hat s', \hat c, \hat c'\}_{m+n_{\rm R}}\big|\mathcal{D}(\mu^2)\big|\{p, f, s, s', c, c'\}_m\big) \\
&\qquad \times \big(\{p, f, s, s', c, c'\}_m\big|\rho_{\rm hard}(\mu^2)\big)
\end{aligned}
$$

Infrared sensitive operator

We can consider a more constructive approach to build the full infrared sensitive operator. This operator basically represents the QCD density operator of a $m \to X$ *(anything)* process.

$$\mathcal{D}(\mu^2) = 1 + \sum_{n=1}^{k} \left[\frac{\alpha_s(\mu^2)}{2\pi}\right]^n \sum_{\substack{n_R=0 \\ n_R+n_V=n}}^{n} \sum_{n_V=0}^{n} \mathcal{D}^{(n_R,n_V)}(\mu^2)$$



$$\left[\frac{\alpha_s(\mu^2)}{2\pi}\right]^2 \mathcal{D}^{(1,1)}(\mu^2)$$

The structure is rather straightforward:

$$\begin{aligned}
&\big(\{\hat{p},\hat{f},\hat{s}',\hat{c}',\hat{s},\hat{c}\}_{m+n_R}\big|\mathcal{D}^{(n_R,n_V)}(\mu^2,\mu_S^2)\big|\{p,f,s',c',s,c\}_m\big) \\
&= \sum_{G\in\text{Graphs}} \int d^d\{\ell\}_{n_V}\ {}_D\big\langle\{\hat{s},\hat{c}\}_{m+n_R}\big|\boldsymbol{V}_L(G;\{\hat{p},\hat{f}\}_{m+n_R},\{\ell\}_{n_V},\mu^2)\big|\{s,c\}_m\big\rangle \\
&\qquad \times \big\langle\{s,c\}_m\big|\boldsymbol{V}_R^\dagger(G;\{\hat{p},\hat{f}\}_{m+n_R},\{\ell\}_{n_V},\mu^2)\big|\{\hat{s},\hat{c}\}_{m+n_R}\big\rangle_D \\
&\quad \times \sum_{I\in\text{Regions(G)}} \big(\{\hat{p},\hat{f}\}_{m+n_R}\big|\mathcal{P}_G(I)\big|\{p,f\}_m\big)\ \underbrace{\Theta_G(I;\{\hat{p},\hat{f}\}_{m+n_R},\{\ell\}_{n_V};\mu_S^2)}
\end{aligned}$$

Constrains the off-shellness of the hard partons

Infrared sensitive operator

- We have to introduce an **ultraviolet cutoff to capture only the IR part** of the amplitudes. At first order level in the real graphs it is just a cut on an infrared sensitive variable of the splitting:

$$\Theta_G(I; \{\hat{p}, \hat{f}\}_{m+n_{\rm R}}, \{\ell\}_{n_{\rm V}}; \mu_{\rm S}^2) \sim \theta(k_\perp^2 < \mu_{\rm S}^2)$$

- The D operator depends on two scales (renormalization scale μ and the **shower scale** $\mu_{\rm S}$) but here we simply set them equal.

$$\mu_{\rm S}^2 = \mu^2$$

- We don't do eikonal approximation in the soft gluon exchange between two external lines because that messes up the **Glauber region**.

- We also need a **momentum mapping**. This can be tricky at higher order level and not necessary the simpler is the better. We prefer "global" momentum mapping.

$\mathrm{N^kLO}$ calculations

Subtractions

Singularities cancel each other here

$$\sigma[O_J] = \overbrace{\big(1\big|\mathcal{O}_J\big[\mathcal{F}(\mu^2)\circ\mathcal{Z}_F(\mu^2)\big]\,\mathcal{D}(\mu^2)}^{\text{Singularities cancel each other here}}\,\underbrace{\mathcal{D}^{-1}(\mu^2)\big|\rho(\mu^2)\big)}_{=|\rho_{\mathrm{H}}(\mu^2))}$$

Hard part, finite in d=4 dimension

$$+\,\mathcal{O}(\alpha_{\mathrm{s}}^{k+1}L^{2k+2})$$
$$+\,\mathcal{O}(\Lambda_{QCD}^2/\mu_J^2)$$

- This is a good approximation as long as
$$\mu^2 < \mu_J^2$$
- the D operator doesn't create resolvable partons, thus
$$\mathcal{D}(\mu^2)\mathcal{O}_J \approx \mathcal{O}_J\mathcal{D}(\mu^2)$$
- otherwise we have to deal with large logarithms,
$$L = \log\frac{\mu^2}{\mu_J^2}$$

Usually $\mathcal{D}^{-1}(\mu^2)$ is constructed by hand and $\mathcal{D}(\mu^2)$ is its inverse.

$$\mathcal{D}^{-1}(\mu_{\mathrm{R}}^2)|\rho(\mu_{\mathrm{R}}^2)) = \overbrace{|\rho^{(0)}(\mu_{\mathrm{R}}^2))}^{\text{Born term}} + \frac{\alpha_{\mathrm{s}}(\mu_{\mathrm{R}}^2)}{2\pi}\overbrace{\Big[|\rho^{(1)}(\mu_{\mathrm{R}}^2)) - \mathcal{D}^{(1)}(\mu_{\mathrm{R}}^2)|\rho^{(0)}(\mu_{\mathrm{R}}^2))\Big]}^{\text{NLO contributions}}$$
$$+\left[\frac{\alpha_{\mathrm{s}}(\mu_{\mathrm{R}}^2)}{2\pi}\right]^2\underbrace{\Big\{|\rho^{(2)}(\mu_{\mathrm{R}}^2)) - \mathcal{D}^{(1)}(\mu_{\mathrm{R}}^2)|\rho^{(1)}(\mu_{\mathrm{R}}^2)) - \big[\mathcal{D}^{(2)}(\mu_{\mathrm{R}}^2) - \mathcal{D}^{(1)}(\mu_{\mathrm{R}}^2)\mathcal{D}^{(1)}(\mu_{\mathrm{R}}^2)\big]|\rho^{(0)}(\mu_{\mathrm{R}}^2))\Big\}}_{\text{NNLO contributions}}$$
$$+\,\mathcal{O}(\alpha_{\mathrm{s}}^3)$$

N^kLO calculations

We define an operator that is **finite** and **doesn't** change the number of patrons and their momenta and flavours in such way that

$$\left(1\right| \underbrace{\mathcal{V}(\mu^2)} = \left(1\right| \underbrace{\left[\mathcal{F}(\mu^2) \circ \mathcal{Z}_F(\mu^2)\right] \mathcal{D}(\mu^2) \mathcal{F}^{-1}(\mu^2)}$$

- IR **finite** operator
- **doesn't** create new patrons
- **doesn't** change momenta or flavours
- its definition is **ambiguous**

- IR **singular** operator
- **does** create new patrons
- **does** change momenta and flavours

With the help of this we can define a normalised IR singular operator as

$$\mathcal{X}_1(\mu^2) = \left[\mathcal{F}(\mu^2) \circ \mathcal{Z}_F(\mu^2)\right] \mathcal{D}(\mu^2) \mathcal{F}^{-1}(\mu^2) \mathcal{V}^{-1}(\mu^2)$$

$\Rightarrow$ $\left(1\right| \mathcal{X}_1(\mu^2) = \left(1\right|$

from definition

The cross section can be written as

$$\sigma[O_J] = \left(1\right| \underbrace{\mathcal{O}_J\, \mathcal{X}_1(\mu^2)}\, \mathcal{V}(\mu^2) \mathcal{F}(\mu^2) \left|\rho_{\mathrm{H}}(\mu^2)\right)$$

when we don't have to worry about large logs, these operators **commute**

Useful notations

It is proven to be useful to generalise the procedure of defining operator $\mathcal{V}(\mu^2)$ from $\mathcal{D}(\mu^2)$.

Let A be a linear operator in the statistical space (may or mayn't change the number of partons):

$$\mathcal{A}\big|\{p,f,c,c',s,s'\}_m\big) = \int d\{\hat{p},\hat{f},\hat{c},\hat{c}',\hat{s},\hat{s}'\}_{\hat{m}}\,\big|\{\hat{p},\hat{f},\hat{c},\hat{c}',\hat{s},\hat{s}'\}_{\hat{m}}\big)\,\big(\{\hat{p},\hat{f},\hat{c},\hat{c}',\hat{s},\hat{s}'\}_{\hat{m}}\big|\mathcal{A}\big|\{p,f,c,c',s,s'\}_m\big)$$

We define a mapping, $[\cdot]_{\mathbb{P}} : \mathcal{A} \longrightarrow \mathcal{B} = [\mathcal{A}]_{\mathbb{P}}$, in such a way that

$$\mathcal{B}\big|\{p,f,c,c',s,s'\}_m\big) = \int d\{\hat{c},\hat{c}',\hat{s},\hat{s}'\}_m\,\big|\{p,f,\hat{c},\hat{c}',\hat{s},\hat{s}'\}_m\big)\,\big(\{p,f,\hat{c},\hat{c}',\hat{s},\hat{s}'\}_m\big|\mathcal{B}\big|\{p,f,c,c',s,s'\}_m\big)$$

and

$$\big(1\big|[\mathcal{A}]_{\mathbb{P}} = \big(1\big|\mathcal{A}$$

The combination $\mathcal{A} - [\mathcal{A}]_{\mathbb{P}}$ appears frequently, thus it is useful to define: $[\mathcal{A}]_{1-\mathbb{P}} = \mathcal{A} - [\mathcal{A}]_{\mathbb{P}}$.

$$\mathcal{V}(\mu_{\mathrm{R}}^2) = \Big[\big[\mathcal{F}(\mu_{\mathrm{R}}^2) \circ \mathcal{Z}_F(\mu_{\mathrm{R}}^2)\big]\,\mathcal{D}(\mu_{\mathrm{R}}^2)\Big]_{\mathbb{P}}\mathcal{F}^{-1}(\mu_{\mathrm{R}}^2)$$

Fixed order cross sections

Parton showers

(only in two slides)

Shower Cross Section

The fixed order cross section is fine as long as we can calculate at "**all order level**". But life is not that easy…

- **truncated** at NLO, NNLO level
- prefers **large scale,** $\mu^2 \approx Q^2$

$$\sigma[O_J] = \big(1\big|\underbrace{\mathcal{O}_J\,\mathcal{X}_1(\mu^2)}\,\mathcal{V}(\mu^2)\overbrace{\mathcal{F}(\mu^2)\,\big|\rho_{\rm H}(\mu^2)\big)}$$

- prefers **small scale**, $\mu^2 \ll \mu_J^2$
- that is **in conflict** with the hard part

- Choose a **hard scale**, $\mu_{\rm H}^2 \approx Q^2$
- Choose a **cutoff scale**, $\mu_J^2 \gg \mu_{\rm f}^2 \approx 1\mathrm{GeV}^2$
- Insert a unit operator between the hard and "soft"part as,

$$1 = \mathcal{X}_1^{-1}(\mu_{\rm f}^2)\mathcal{X}_1(\mu_{\rm f}^2)$$

$$\sigma[O_J] = \underbrace{\big(1\big|\mathcal{O}_J\,\mathcal{X}_1(\mu_{\rm f}^2)}_{=(1|\mathcal{O}_J}\,\overbrace{\mathcal{X}_1^{-1}(\mu_{\rm f}^2)\mathcal{X}_1(\mu_{\rm H}^2)}^{\mathcal{U}(\mu_{\rm f}^2,\mu_{\rm H}^2)}\,\mathcal{V}(\mu_{\rm H}^2)\mathcal{F}(\mu_{\rm H}^2)\,\big|\rho_{\rm H}(\mu_{\rm H}^2)\big)$$

No resolvable radiation come from $\mathcal{X}_1(\mu_{\rm f}^2)$ operator, thus these operators commute,

$$\mathcal{O}_J\mathcal{X}_1(\mu_{\rm f}^2) \approx \mathcal{X}_1(\mu_{\rm f}^2)\mathcal{O}_J.$$

$$\mathcal{U}(\mu_{\rm f}^2,\mu_{\rm H}^2) = \mathbb{T}\exp\left\{\int_{\mu_{\rm f}^2}^{\mu_{\rm H}^2}\frac{d\mu^2}{\mu^2}\,\mathcal{S}(\mu^2)\right\}$$

$$\frac{1}{\mu^2}\mathcal{S}(\mu^2) = \lim_{\epsilon\to 0}\mathcal{X}_1^{-1}(\mu^2)\frac{d\mathcal{X}_1(\mu^2)}{d\mu^2}$$

First order shower

The generators of the unitary shower can be expanded in the coupling:

$$S(\mu^2) = \frac{\alpha_s(\mu^2)}{2\pi} S^{(1)}(\mu^2) + \left[\frac{\alpha_s(\mu^2)}{2\pi}\right]^2 S^{(2)}(\mu^2) + \cdots$$

and the first order term is rather simple

$$\frac{1}{\mu^2} S^{(1)}(\mu^2) = \left[\underbrace{\mathcal{F}(\mu_R^2) \frac{\partial \mathcal{D}^{(1,0)}(\mu^2, \mu_S^2)}{\partial \mu_S^2} \mathcal{F}^{-1}(\mu^2)} - \underbrace{\frac{\partial \left[\mathcal{F}(\mu_R^2) \mathcal{D}^{(1,0)}(\mu^2, \mu_S^2)\right]_{\mathbb{P}}}{\partial \mu_S^2} \mathcal{F}^{-1}(\mu^2)} + \underbrace{\mathrm{Im} \frac{\partial \mathcal{D}^{(0,1)}(\mu^2, \mu_S^2)}{\partial \mu_S^2}} \right]_{\mu_S^2 = \mu^2}$$

Real operator
all the quantum numbers of the emitted parton is **resolved**

Integrated real operator
- all the quantum numbers of the emitted parton is **integrated out**
- it is **not** the contribution of the virtual graphs

Glauber gluon
imaginary part of the virtual graphs

$\sim \mathrm{i}\pi$

Note, the first order kernel is independent of the real part of the virtual graphs.

Fixed order cross sections

Parton showers

Summing large logarithms with parton showers

Summing logarithms

I don't trust in eye measure to clam LL or NLL accuracy of any parton shower. The only way to check the summation property of the shower is to **gain analytical control** on the shower cross section. Is it *possible* to do it? Is it *simple*?

$$\sigma[O_J] = \big(1\big|\mathcal{O}_J\,\mathbb{T}\exp\left\{\int_{\mu_{\rm f}^2}^{\mu_{\rm H}^2}\frac{d\mu^2}{\mu^2}\,\mathcal{S}(\mu^2)\right\}\mathcal{V}(\mu_{\rm H}^2)\mathcal{F}(\mu_{\rm H}^2)\,\big|\rho_{\rm H}(\mu_{\rm H}^2)\big)$$

- **"infinite"** number of partons
- make measurement on these multi-parton states
- **impossible task** to study the log structure analytically

We should **reformulate** the shower cross section, in such a way that:

➠ more suitable for analytical studies

➠ **without** extra approximation (all the approximations have been done in the shower operator S)

➠ the effect of the measurement operator should be **exponentiated**

We want to test the log summation property of the parton shower cross algorithms

➠ study observables that exponentiates (**thrust**, Drell-Yan pT-distributions,…)

➠ analytical results are available

Preparing observables

Consider the Dell-Yan kT distribution:

$$\hat{\mathcal{O}}(\boldsymbol{k}_\perp)\big|\{p,f,...\}_m\big) = (2\pi)^2\delta^{(2)}(\boldsymbol{k}_\perp - \boldsymbol{k}_Z(\{p\}_m))\big|\{p,f,...\}_m\big)$$

This operator is not invertible, but its **Fourier transform** is,

$$\mathcal{O}(\boldsymbol{b})\big|\{p,f,...\}_m\big) = e^{\mathrm{i}\boldsymbol{b}\cdot\boldsymbol{k}_Z(\{p\}_m)}\big|\{p,f,...\}_m\big)\ , \qquad \mathcal{O}^{-1}(\boldsymbol{b})\big|\{p,f,...\}_m\big) = e^{-\mathrm{i}\boldsymbol{b}\cdot\boldsymbol{k}_Z(\{p\}_m)}\big|\{p,f,...\}_m\big)\ .$$

Similarly for **thrust**, we use **Laplace transformation** to make the measurement operator invertible,

$$\mathcal{O}(\nu)\big|\{p,f,...\}_m\big) = e^{-\nu\tau(\{p\}_m)}\big|\{p,f,...\}_m\big)\ , \qquad \mathcal{O}^{-1}(\nu)\big|\{p,f,...\}_m\big) = e^{\nu\tau(\{p\}_m)}\big|\{p,f,...\}_m\big)$$

The formalism can deal with measurement operator that has an inverse, thus we almost **always need** some kind of **proxy** to do the analytical studies of the parton showers. Sometimes it is just a simple integral transformation, sometimes a generating functional. It is a good guideline to follow the footsteps of the analytic calculation.

$\hat{\mathcal{O}}(\boldsymbol{v}) \Longrightarrow \mathcal{O}(\boldsymbol{r})$ and $\mathcal{O}(\boldsymbol{r})$ always **has an inverse** over the **whole** statistical space

Observable dependent shower

We define an operator that is **finite** and **doesn't** change the number of patrons and their momenta and flavours but this time **with observable dependence**

$$\mathcal{Y}(\mu^2;\boldsymbol{r}) = \big[\mathcal{O}(\boldsymbol{r})\big[\mathcal{F}(\mu^2)\circ\mathcal{Z}_F(\mu^2)\big]\,\mathcal{D}(\mu^2)\mathcal{O}^{-1}(\boldsymbol{r})\big]_{\mathbb{P}} \times\Big(\big[\big[\mathcal{F}(\mu^2)\circ\mathcal{Z}_F(\mu^2)\big]\,\mathcal{D}(\mu^2)\big]_{\mathbb{P}}\Big)^{-1}$$

- IR **finite** operator
- **doesn't** create new patrons
- **doesn't** change momenta or flavours
- its definition obviously is **ambiguous**
- **normalised**

$$\mathcal{O}(\boldsymbol{r}) = 1 \implies \mathcal{Y}(\mu^2;\boldsymbol{r}) = 1$$

From the definition, it is easy to show that

$$\big(1\big|\mathcal{Y}(\mu^2;\boldsymbol{r})\mathcal{O}(\boldsymbol{r}) = \big(1\big|\mathcal{O}(\boldsymbol{r})\,\mathcal{U}(\mu_{\rm f}^2,\mu^2)$$

and the shower cross section becomes

measurement **after the shower** (many patrons)

$$\sigma(\boldsymbol{r}) = \big(1\big|\mathcal{Y}(\mu_{\rm H}^2,\boldsymbol{r})\,\mathcal{V}(\mu_{\rm H}^2)\mathcal{O}(\boldsymbol{r})\,\mathcal{F}(\mu_{\rm H}^2)\,\big|\rho_{\rm H}(\mu_{\rm H}^2)\big) = \big(1\big|\mathcal{O}(\boldsymbol{r})\,\mathcal{U}(\mu_{\rm f}^2,\mu_{\rm H}^2)\,\mathcal{V}(\mu_{\rm H}^2)\,\mathcal{F}(\mu_{\rm H}^2)\,\big|\rho_{\rm H}(\mu_{\rm H}^2)\big)$$

measurement on the **hard state** (only few patrons)

It is **really** an **equal sign**!

Observable dependent shower

The $\mathcal{Y}(\mu^2;\boldsymbol{r})$ operator can be **exponentiated** in the usual way,

$$\mathcal{Y}(\mu_{\rm H}^2;\boldsymbol{r}) = \mathbb{T}\exp\left\{\int_{\mu_{\rm f}^2}^{\mu_{\rm H}^2}\frac{d\mu^2}{\mu^2}\mathcal{S}_{\mathcal{Y}}(\mu^2;\boldsymbol{r})\right\} \quad , \quad \text{with} \quad \mathcal{Y}(\mu_{\rm f}^2;\boldsymbol{r}) = 1$$

where

$$\frac{1}{\mu^2}\mathcal{S}_{\mathcal{Y}}(\mu^2;\boldsymbol{r}) = \mathcal{Y}^{-1}(\mu^2;\boldsymbol{r})\frac{d\mathcal{Y}(\mu^2;\boldsymbol{r})}{d\mu^2} \quad .$$

- Here the exponent **has to be an all order** expression to maintain the equality with the shower cross section.
- The operator $\mathcal{S}_{\mathcal{Y}}(\mu^2;\boldsymbol{r})$ contains large logarithms of $L(\boldsymbol{r})$.
- We can relate the $\mathcal{S}_{\mathcal{Y}}(\mu^2;\boldsymbol{r})$ operator to the generator of the parton shower $\mathcal{S}(\mu^2)$ via

$$\big(1\big|\mathcal{Y}(\mu^2;\boldsymbol{r})\,\mathcal{S}_{\mathcal{Y}}(\mu^2;\boldsymbol{r}) = \big(1\big|\mathcal{Y}(\mu^2;\boldsymbol{r})\,\mathcal{O}(\boldsymbol{r})\,\mathcal{S}(\mu^2)\,\mathcal{O}^{-1}(\boldsymbol{r})$$

with the help of the $[\cdot]_{\mathbb{P}}$ operation we can extract $\mathcal{S}_{\mathcal{Y}}(\mu^2;\boldsymbol{r})$ as

$$\mathcal{S}_{\mathcal{Y}}(\mu^2;\boldsymbol{r}) = \left[\mathcal{Y}(\mu^2;\boldsymbol{r})\,\mathcal{O}(\boldsymbol{r})\,\mathcal{S}(\mu^2)\,\mathcal{O}^{-1}(\boldsymbol{r})\right]_{\mathbb{P}} - \left[\mathcal{Y}(\mu^2;\boldsymbol{r}) - 1\right]\mathcal{S}_{\mathcal{Y}}(\mu^2;\boldsymbol{r})$$

and this can be solved **recursively** order by order (in powers of the shower generator $\mathcal{S}(\mu^2)$).

Observable dependent shower

We have **two equations and two unknowns**, so we can solve them recursively:

$$\mathcal{S}_{\mathcal{Y}}(\mu^2;\nu) = \sum_{k=1}^{\infty} \mathcal{S}_{\mathcal{Y}}^{[k]}(\mu^2;\nu)$$

$$\mathcal{Y}(\mu^2;\nu) = 1 + \sum_{k=1}^{\infty} \mathcal{Y}^{[k]}(\mu^2;\nu)$$

At **first order** level we have

$$\mathcal{S}_{\mathcal{Y}}^{[1]}(\mu^2;\boldsymbol{r}) = \big[\mathcal{O}(\boldsymbol{r})\,\mathcal{S}(\mu^2)\,\mathcal{O}^{-1}(\boldsymbol{r})\big]_{\mathbb{P}}$$

$$\mathcal{Y}^{[1]}(\mu^2;\boldsymbol{r}) = \int_{\mu_{\rm f}^2}^{\mu^2} \frac{d\bar{\mu}^2}{\bar{\mu}^2}\big[\mathcal{O}(\boldsymbol{r})\,\mathcal{S}(\bar{\mu}^2)\,\mathcal{O}^{-1}(\boldsymbol{r})\big]_{\mathbb{P}}$$

The **second order** generator is a little bit more complicated:

$$\mathcal{S}_{\mathcal{Y}}^{[2]}(\mu^2;\boldsymbol{r}) = \int_{\mu_{\rm f}^2}^{\mu^2} \frac{d\bar{\mu}^2}{\bar{\mu}^2}\Big[\big[\mathcal{O}(\boldsymbol{r})\,\mathcal{S}(\bar{\mu}^2)\,\mathcal{O}^{-1}(\boldsymbol{r})\big]_{\mathbb{P}}\big[\mathcal{O}(\boldsymbol{r})\,\mathcal{S}(\mu^2)\,\mathcal{O}^{-1}(\boldsymbol{r})\big]_{1-\mathbb{P}}\Big]_{\mathbb{P}}$$

Now the shower cross section (in a kind of analytical form) is

$$\sigma(\boldsymbol{r}) = \big(1\big|\mathbb{T}\exp\left\{\int_{\mu_{\rm f}^2}^{\mu_{\rm H}^2}\frac{d\mu^2}{\mu^2}\left(\mathcal{S}_{\mathcal{Y}}^{[1]}(\mu^2,\boldsymbol{r}) + \sum_{k=2}^{\infty}\mathcal{S}_{\mathcal{Y}}^{[k]}(\mu^2,\boldsymbol{r})\right)\right\}\mathcal{V}(\mu_{\rm H}^2)\mathcal{O}(\boldsymbol{r})\,\mathcal{F}(\mu_{\rm H}^2)\,\big|\rho_{\rm H}(\mu_{\rm H}^2)\big)$$

For many observables the **exponentiated single emission** operator provides NLL accuracy.

This should lead to only **subleading** logs (NNLL,...).

Thrust in e+e- annihilation

In this case the hard process at Born level is very simple, it is proportional to a single basis vector only with a quark-antiquark pair:

$$\big|\rho_{\rm H}(\mu_{\rm H}^2)\big) \propto \big|\{p,f,c,c\}_2\big)$$

This is always eigenvector of the exponent, thus the exponentiation is trivial:

$$\mathcal{S}_{\mathcal{Y}}^{[k]}(\mu^2;\nu)\big|\{p,f,c,c\}_2\big) = \lambda_{\mathcal{Y}}^{[k]}(\mu^2/Q^2;\nu)\big|\{p,f,c,c\}_2\big)$$

With this the cross section is rather simple,

This is the "**golden nugget**". Any reasonable parton shower algorithm agrees with the analytic result.

$$\frac{\sigma(\boldsymbol{r})}{\sigma_0} = \exp\left\{\int_0^1 \frac{dx}{x}\left(\overbrace{\lambda_{\mathcal{Y}}^{[1]}(x,\nu)} + \underbrace{\sum_{k=2}^{\infty}\lambda_{\mathcal{Y}}^{[k]}(x,\nu)}\right)\right\} + \cdots$$

This is the shower generated "**junk**". This has to be subleading log contribution.

- We can study analytically the exponent when it is possible,

$$I^{[k]}(\nu) = \int_0^1 \frac{dx}{x}\lambda_{\mathcal{Y}}^{[k]}(x,\nu)$$

- When it is hard to test analytically, we can calculate the exponent numerically and test its large log behaviour in term of $\log(\nu)$.

$$I^{[k]}(\nu) = \sum_{n=k}^{\infty}\left[\frac{\alpha_{\rm s}(Q^2/\nu)}{2\pi}\right]^n I_n^{[k]}(\nu)$$

- For NLL accuracy we should have

$$I_n^{[k]}(\nu) \sim \log^{n-1}(\nu)$$

for every $k > 1$.

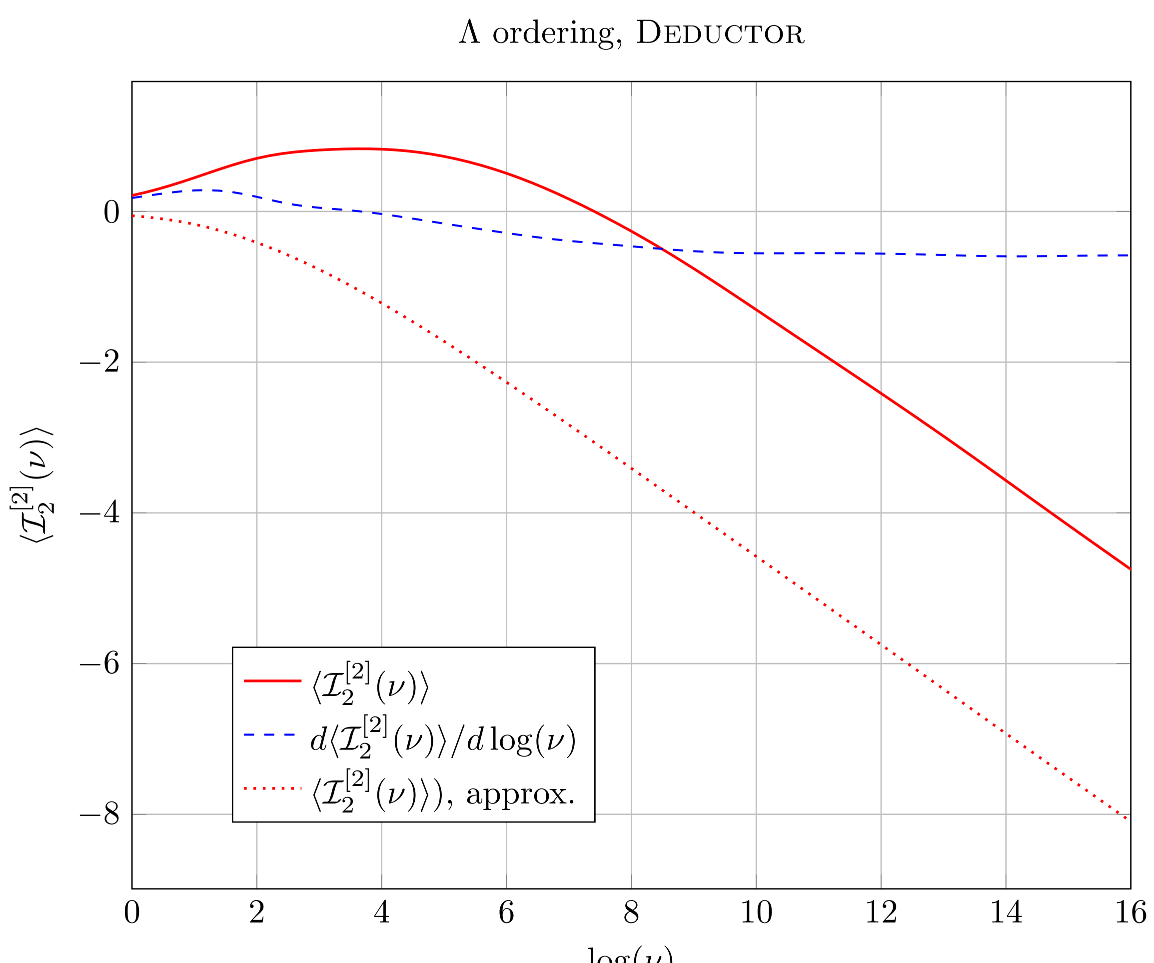
Thrust in e+e- annihilation

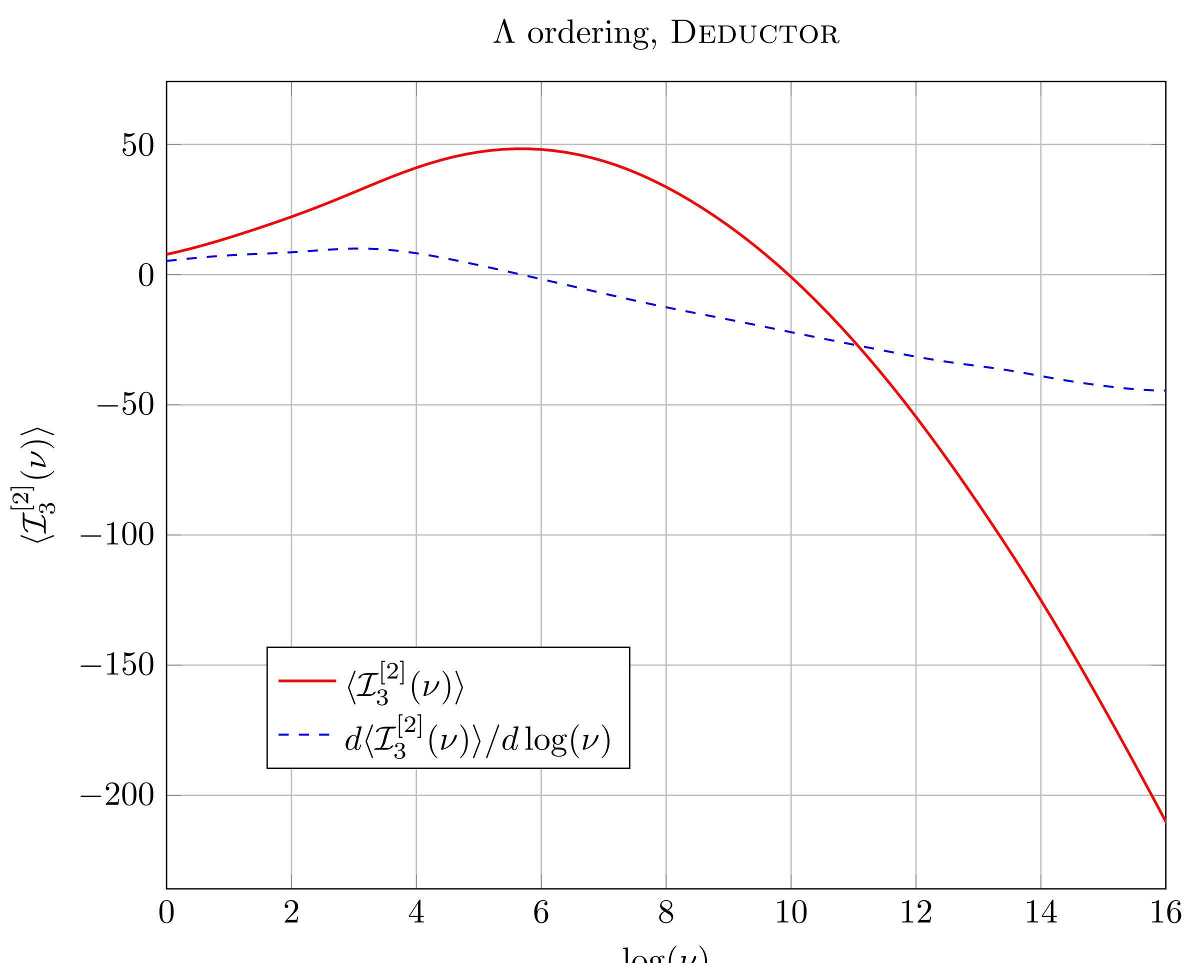
DEDUCTOR Lambda ordered shower

- The ordering variable is the virtuality divided by the mother parton energy

$$\Lambda^2 = \frac{(\hat{p}_l + \hat{p}_{m+1})^2}{2p_l \cdot Q} Q^2$$

- **Global** momentum mapping
- Proper soft gluon treatment with **full colour** evolution
- In this case we **can proof analytically** that the shower sums up large logarithms at NLL level

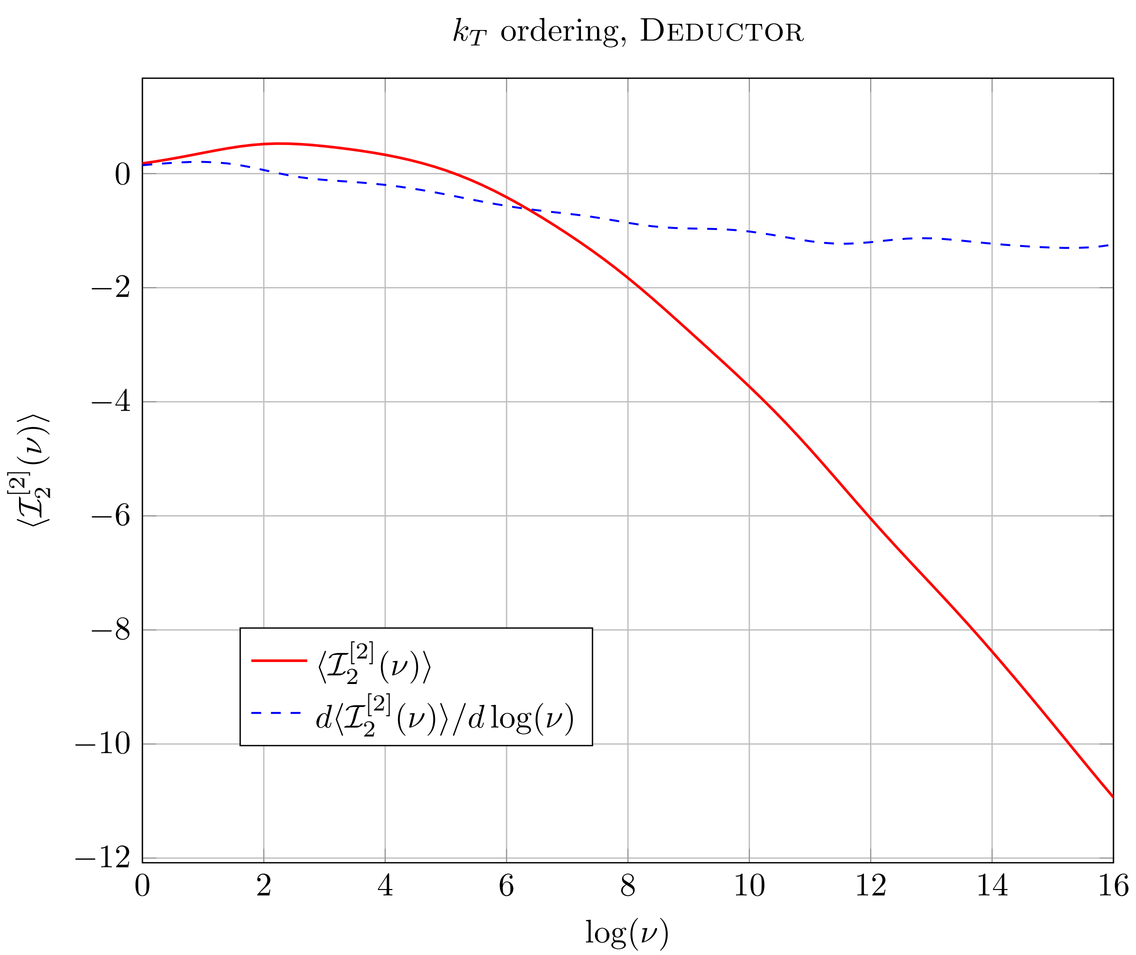


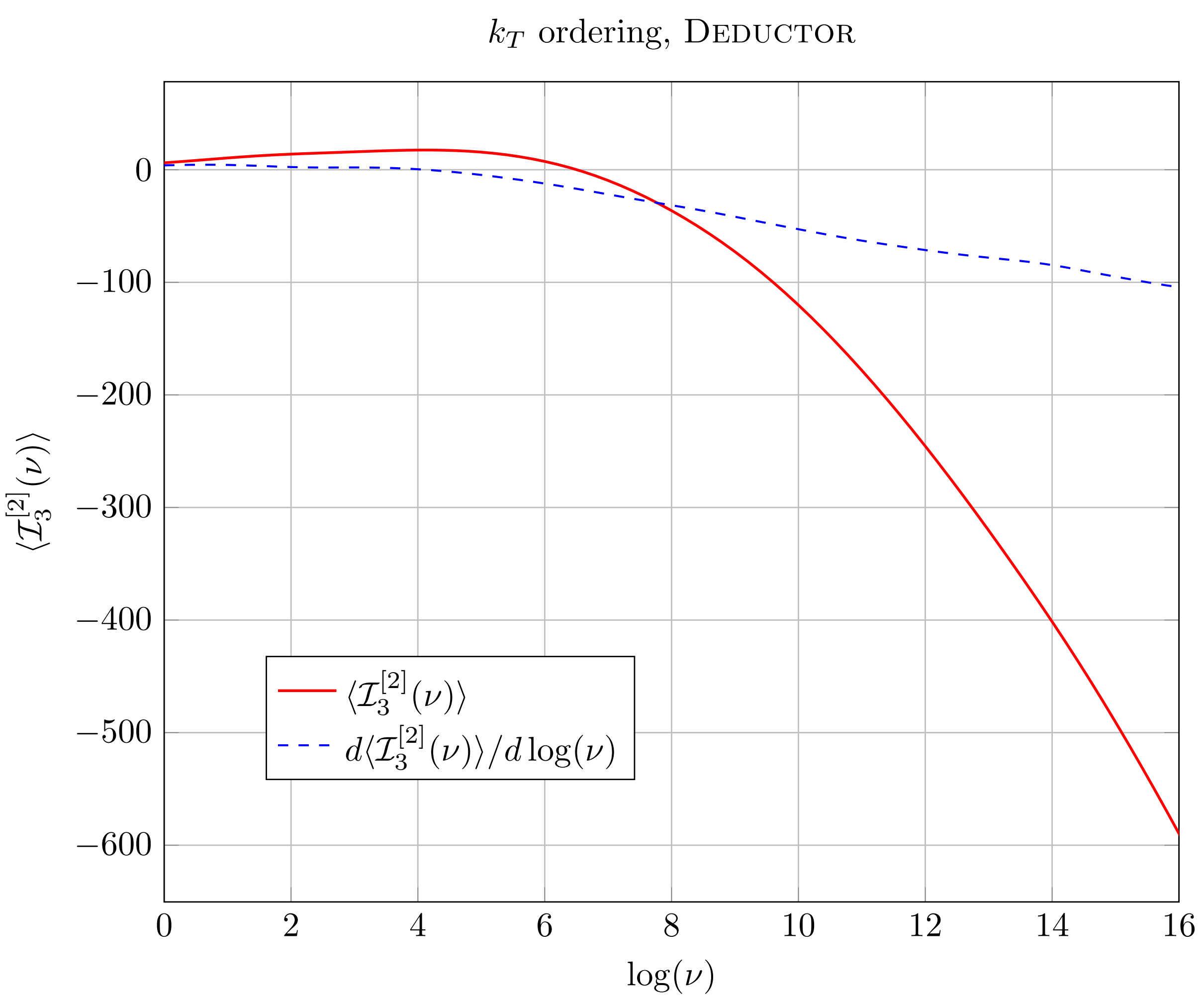


Thrust in e+e- annihilation

DEDUCTOR kT ordered shower

- The ordering variable is the **transverse momentum** of the splitting
- **Global** momentum mapping
- Proper soft gluon treatment with **full colour** evolution
- In this case we **cannot proof analytically** that the shower sums up large logarithms at NLL level
- We check numerically the first couple of $I_n^{[2]}(\nu)$ coefficients.
- It looks OK for k=2 and can be explained by real-virtual cancellation, but hard to see what happens for k > 2.





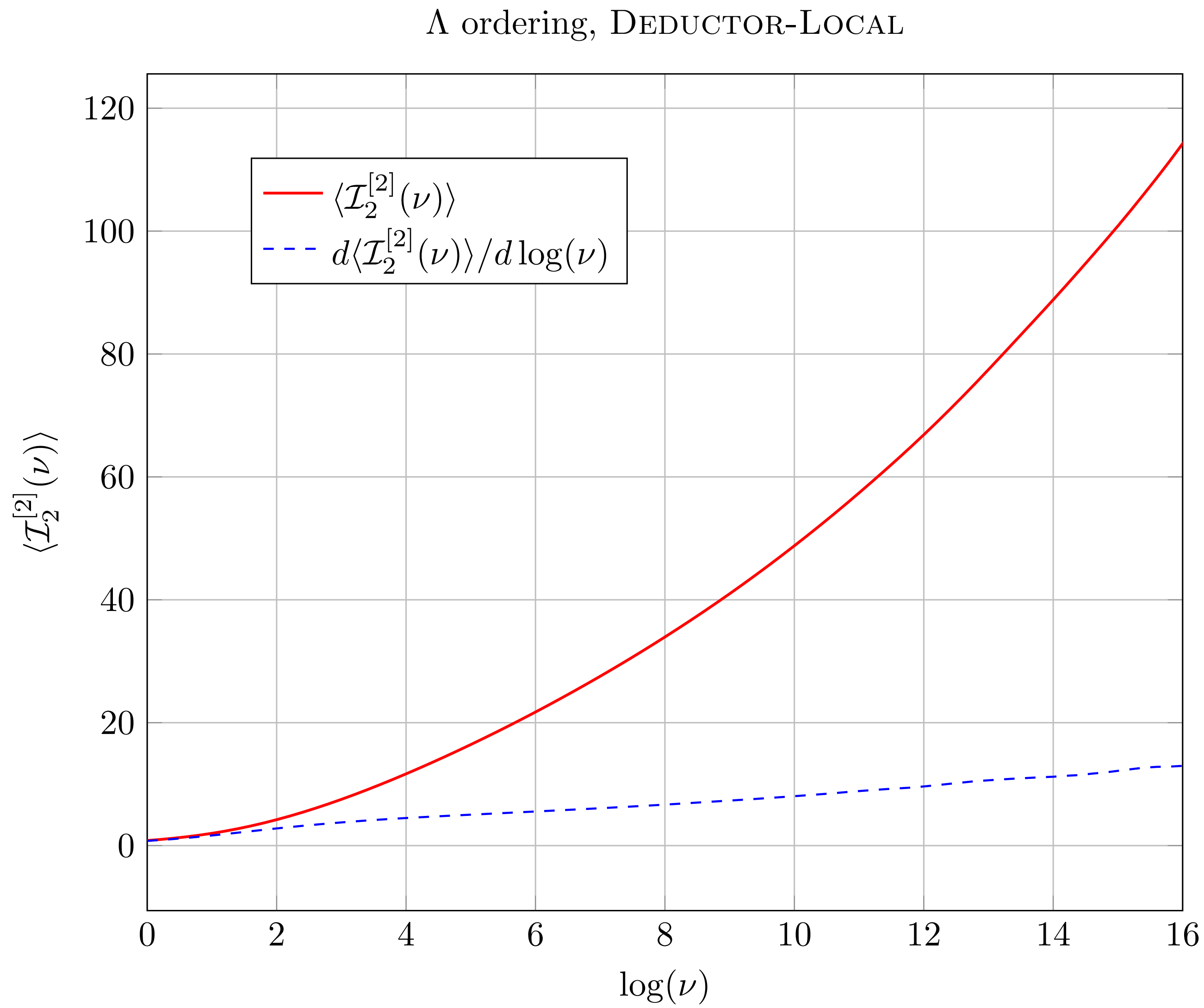
Thrust in e+e- annihilation

DEDUCTOR Lambda ordered shower

- The ordering variable is the virtuality divided by the mother parton energy

$$\Lambda^2 = \frac{(\hat{p}_l + \hat{p}_{m+1})^2}{2p_l \cdot Q} Q^2$$

- **Local** momentum mapping (Catani-Seymour mapping)
- Proper soft gluon treatment with **full colour** evolution
- **Only LL** accuracy can be achieved.

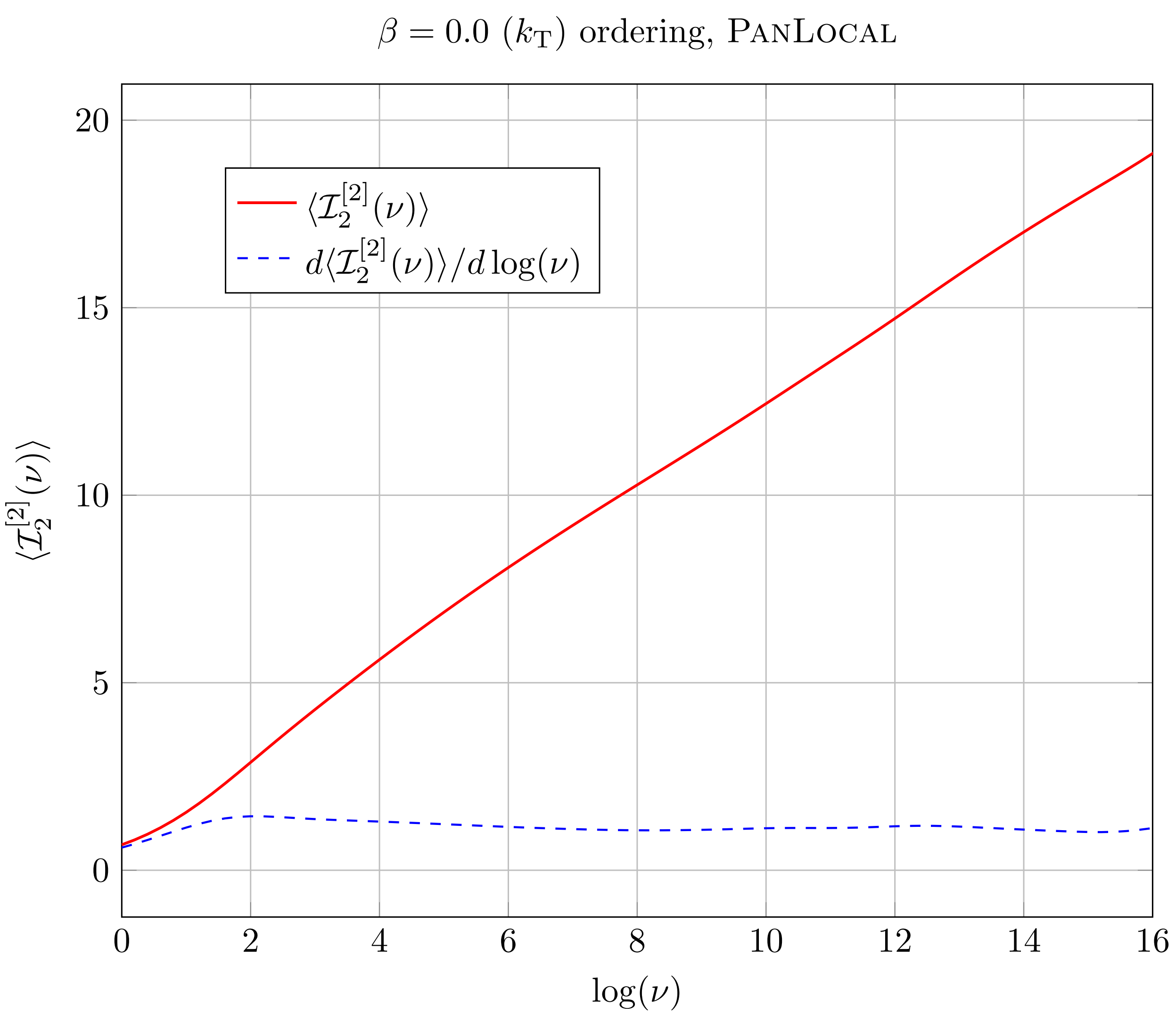


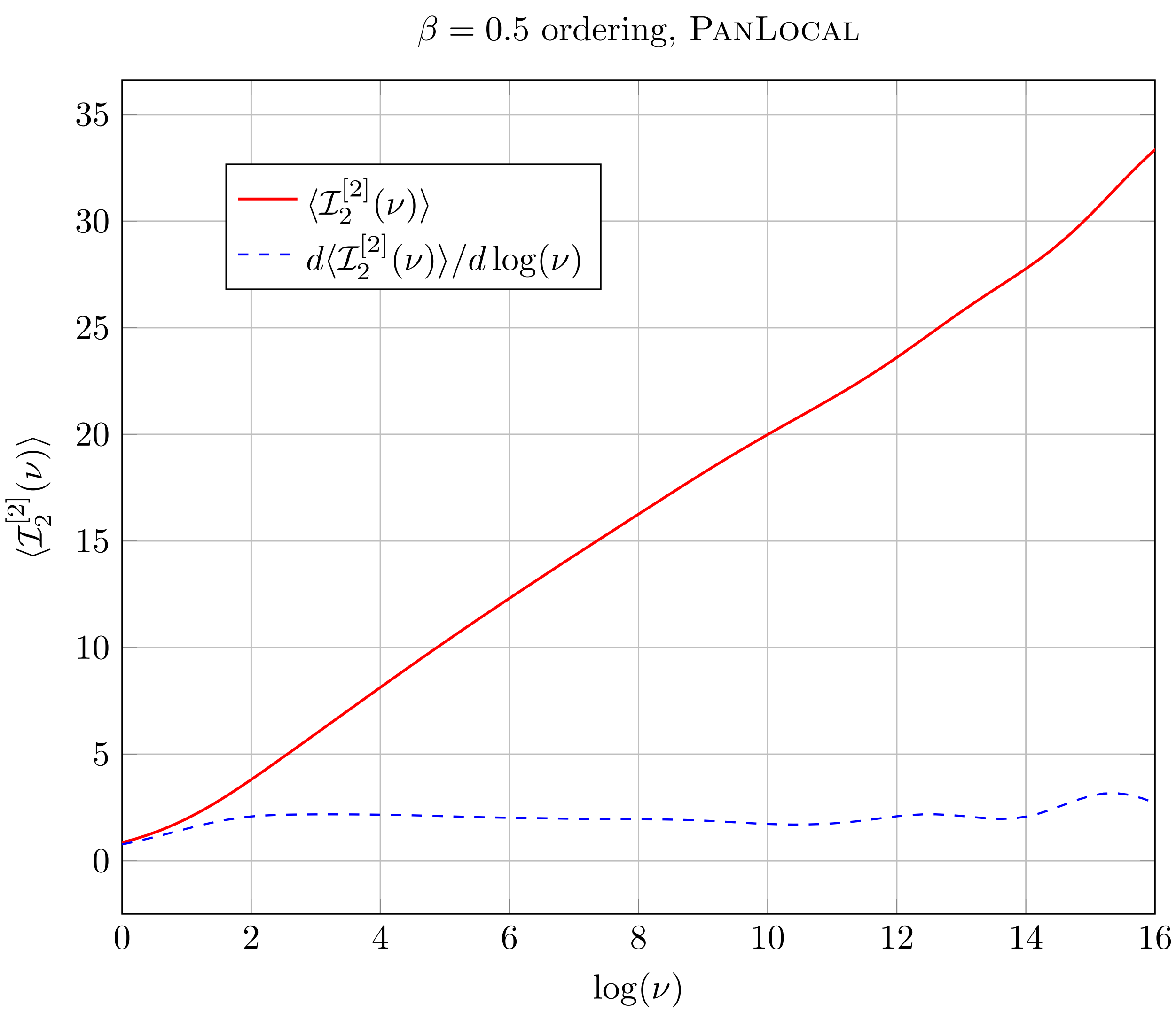
Thrust in e+e- annihilation

Dasgupta *et al.*, ***Phys.Rev.Lett.* 125** (2020) 5, 052002

PanLocal shower $\beta = 0, 0.5$

- The ordering variable is transverse momentum based
- **Local** momentum mapping (it is Catani-Seymour mapping)
- Proper soft gluon treatment with **full colour** evolution (this is not in the original definition)
- It works similarly like the **Deductor** kT ordered shower for $\beta = 0, 0.5$, but fails for $\beta = 1$ (only LL accuracy).





Conclusion

- General and unified scheme for fixed order and parton shower calculation.
- We managed to reformulate the shower cross section in such a way to be able to compare with analytical calculations.
- As long as we do all order calculation, all the three approaches lead to the same cross section.
  - Fixed order calculations are truncated in $\alpha_s(\mu^2)$ at **cross section** level.
  - Parton shower calculations are truncated in $\alpha_s(\mu^2)$ in the **shower exponent.**
  - The "shower resummation formulae" is truncated in $\alpha_s(\mu^2)L$ in the "Sudakov" exponent.
- We extensively studied the thrust distribution in e+e- annihilation.
  - We were able to prove analytically the NLL summation property only in lambda ordered **DEDUCTOR**.
  - With other shower schemes we were showed numerically that $I^{[2]}(\nu)$ is only subleading log contribution. We could not say anything about the higher order contributions.

Outlook

- We want to test more observables
  - Jet rates in e+e- annihilation
  - Drell-Yan kT distribution with and without threshold logarithm
  - …
- Our shower scheme is still **not general enough**. It cannot accommodate the **angular ordered shower** correctly and systematically.
- In the recent years there have been lots of progress on NNLO fixed order calculations. This is a good base to start to think about **NLO parton shower**… *It will be painful…*