QCD Power Corrections

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Trento Workshop

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- Phenomenology of power corrections
- Asymptotics of perturbation theory
- 'Dispersive' approach
- Selected results
- Conclusions

Some recent (1998/9) theoretical references:

YuL Dokshitzer, A Lucenti, G Marchesini, GP Salam, hep-ph/9802381, 9812487; VI Zakharov, hep-ph/9802416, 9811294, 9812374; E Stein, M Maul, L Mankiewicz, A Schäfer, hep-ph/9803342; M Dasgupta, GE Smye, BR Webber, hep-ph/9803382; BR Webber, hep-ph/9805484; G Sterman, hep-ph/9806533; GP Korchemsky, hep-ph/9806537; M Beneke, hep-ph/9807443; G Grunberg, hep-ph/9807494; GE Smye, hep-ph/9810292, 9812251; GP Korchemsky, G Sterman, hep-ph/9902341; YuL Dokshitzer, G Marchesini, BR Webber, hep-ph/9905339; E Gardi, G Grunberg, hep-ph/9908458

Phenomenology of power corrections

Many QCD observables differ from the predictions of perturbation theory by terms that vary as inverse (integer) powers of the relevant hard scale Q:

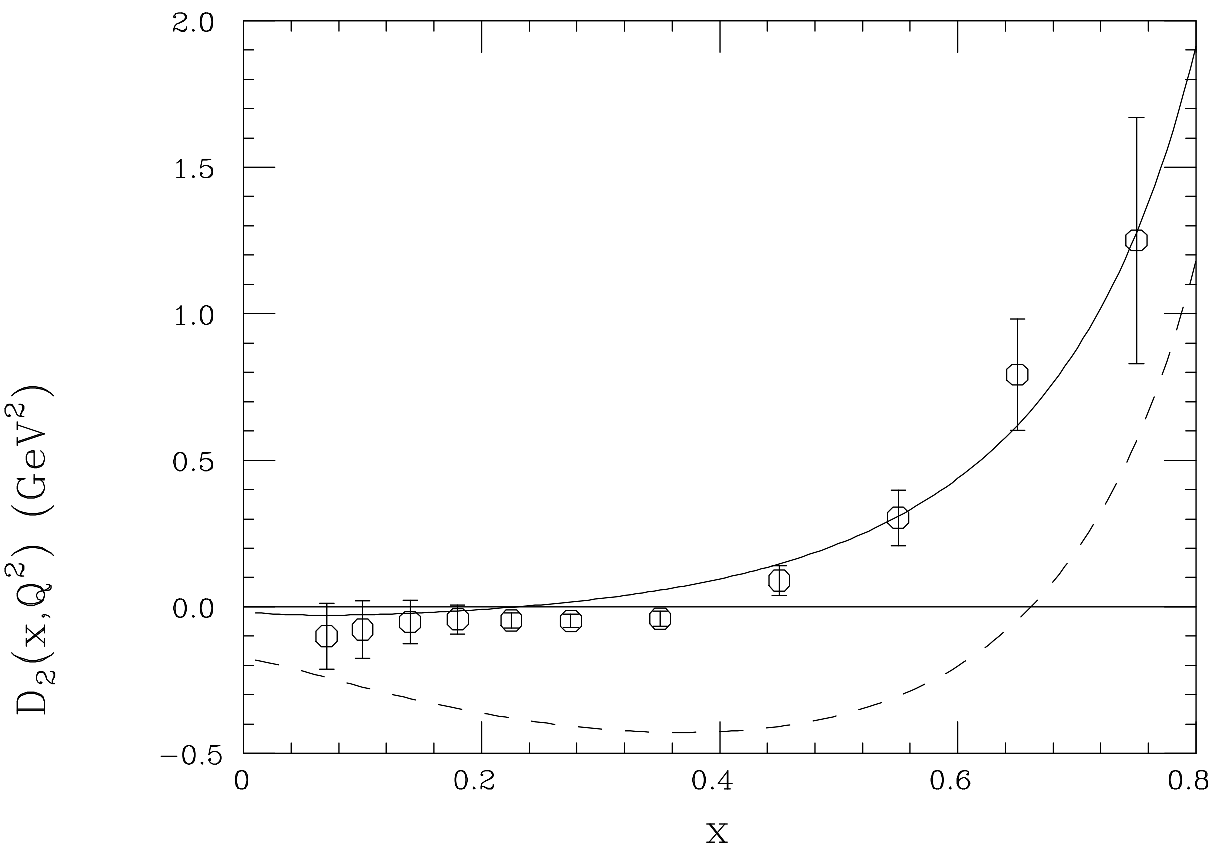
- Deep inelastic structure functions

  Deviations from the perturbative Q^2 dependence generally behave like $1/Q^2$ ("twist 4")

$$F_2(x,Q^2) \simeq F_2^{\rm pert}(x,Q^2)\left[1+\frac{D_2(x,Q^2)}{Q^2}\right]$$

M Dasgupta & BRW, hep-ph/9604388

Data from M Virchaux & A Milsztajn, PLB274(1992)221

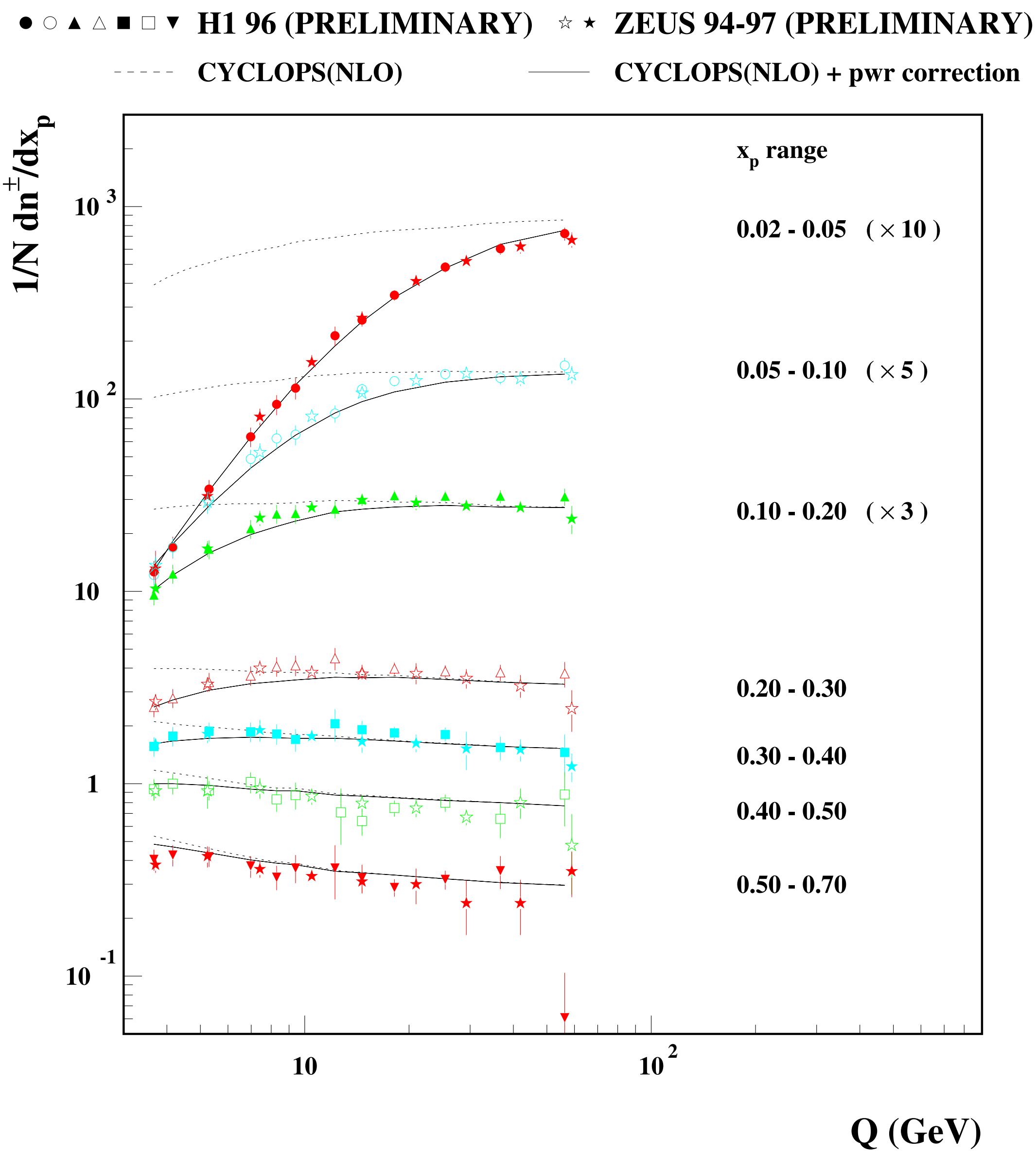


DIS fragmentation functions

The fragmentation function of the current jet (distribution of $x_p = 2|\boldsymbol{p}|/Q$ in the Breit frame) also shows large $1/Q^2$ corrections:

$$F(x_p, Q^2) \simeq F^{\text{pert}}(x_p, Q^2)\left[1 + \frac{0.4\ \text{GeV}^2}{x_p^2\, Q^2}\right]^{-1}$$

P Dixon D Kant & G Thompson, to appear in J Phys G



Event shape variables (e^+e^-)

Event shapes describing the final-state momentum distribution in $e^+e^- \to$ hadrons generally have power corrections that behave like $1/Q$:

Thrust

$$T = \max_{\boldsymbol{n}} \frac{\sum_i |\boldsymbol{p}_i \cdot \boldsymbol{n}|}{\sum_i |\boldsymbol{p}_i|}$$

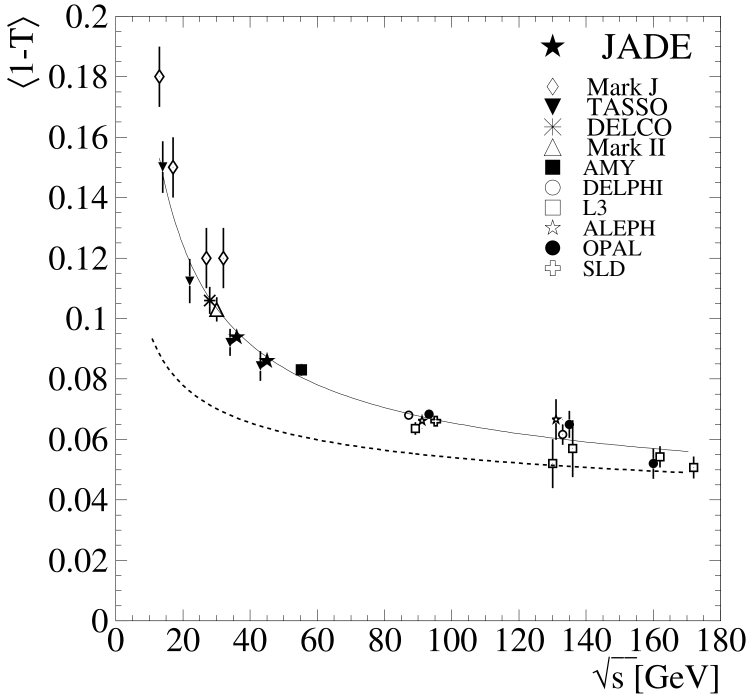
$$\langle 1 - T \rangle = \langle 1 - T \rangle^{\text{pert}} + \delta \langle 1 - T \rangle$$

where

$$\langle 1 - T \rangle^{\text{pert}} = A_T\, \alpha_{\text{S}}(Q^2) + B_T\, \alpha_{\text{S}}^2(Q^2)$$

$$\delta \langle 1 - T \rangle \sim 1 \text{ GeV}/Q$$

JADE Collab, P Movilla Fernandez et al., PITHA 98/21



❖ C-Parameter

$$C = \frac{\sum_{i,j} \left[|\boldsymbol{p}_i||\boldsymbol{p}_j| - (\boldsymbol{p}_i \cdot \boldsymbol{p}_j)^2 / |\boldsymbol{p}_i||\boldsymbol{p}_j| \right]}{\left(\sum_i |\boldsymbol{p}_i|\right)^2}$$

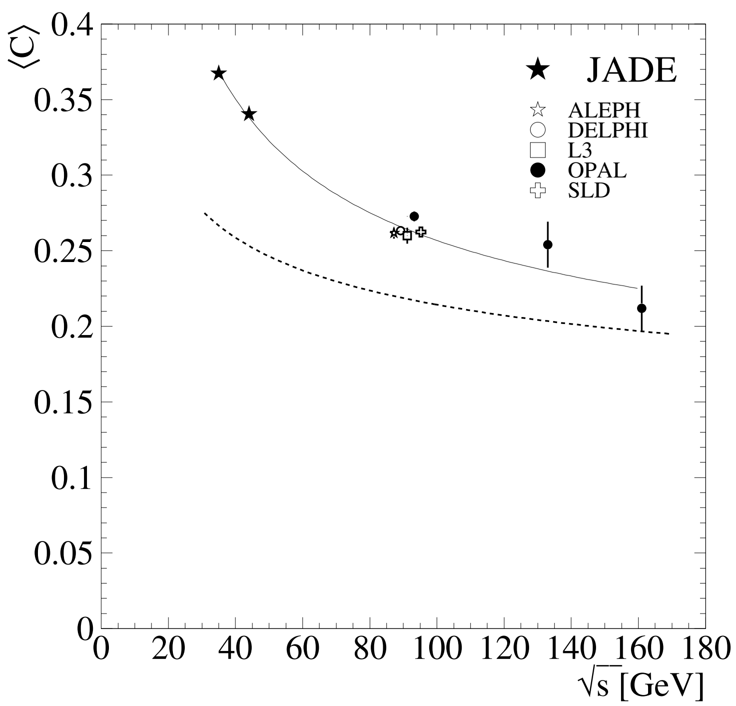
$$\langle C \rangle = \langle C \rangle^{\text{pert}} + \delta \langle C \rangle$$

where

$$\langle C \rangle^{\text{pert}} = A_C \, \alpha_{\text{S}}(Q^2) + B_C \, \alpha_{\text{S}}^2(Q^2)$$

$$\delta \langle C \rangle \sim 4 \text{ GeV}/Q$$

JADE Collab, P Movilla Fernandez et al., PITHA 98/21

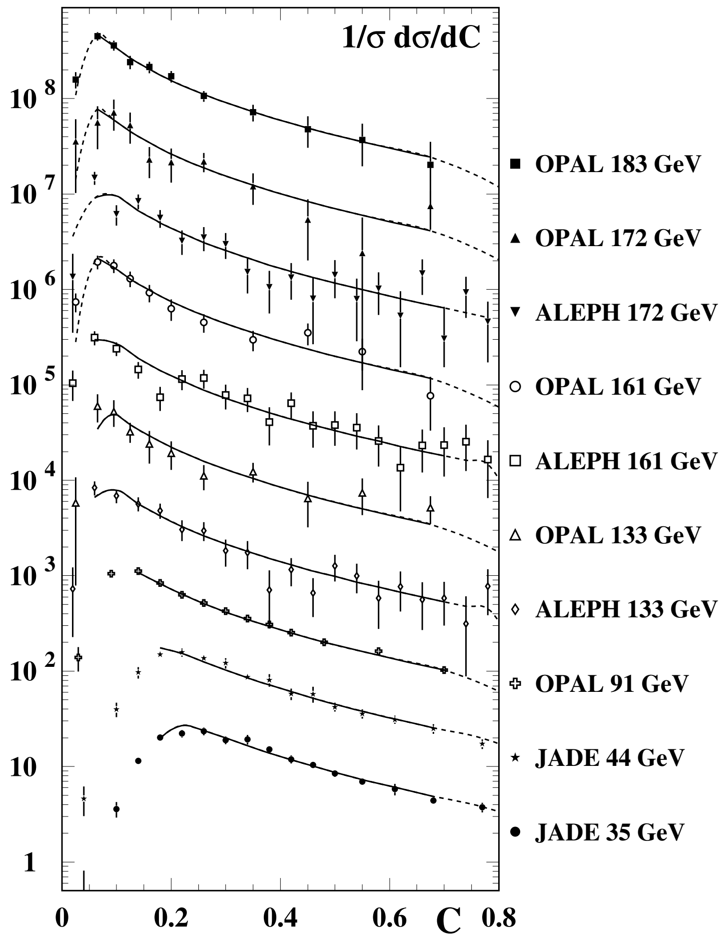


❖ The *distributions* of shape variables have non-perturbative shifts and/or distortions which also behave like $1/Q$:

$$\frac{1}{\sigma}\frac{d\sigma}{dC} \simeq f_C^{\mathrm{pert}}\left(C - \frac{D_C}{Q}\right)$$

where $D_C \simeq 4$ GeV.

P Movilla Fernandez, PITHA 98/24



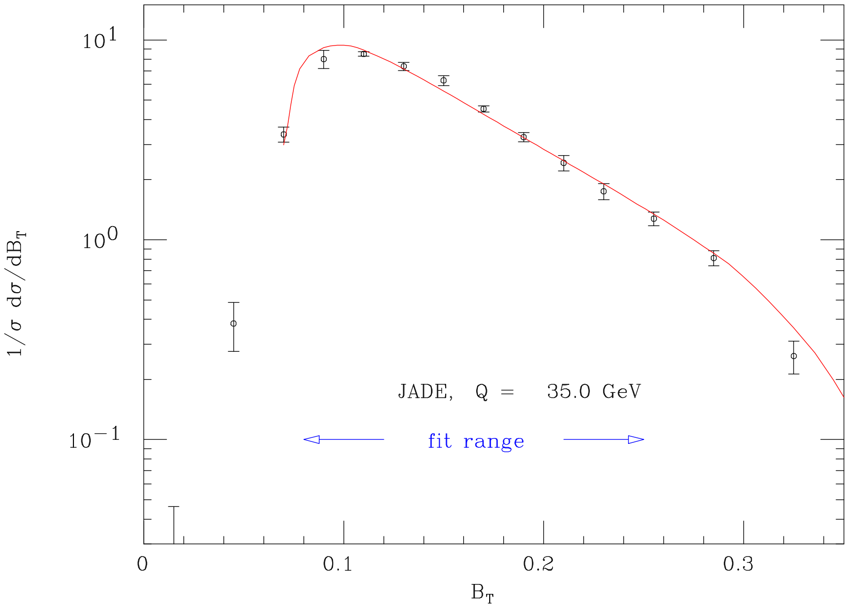
❖ Jet Broadening

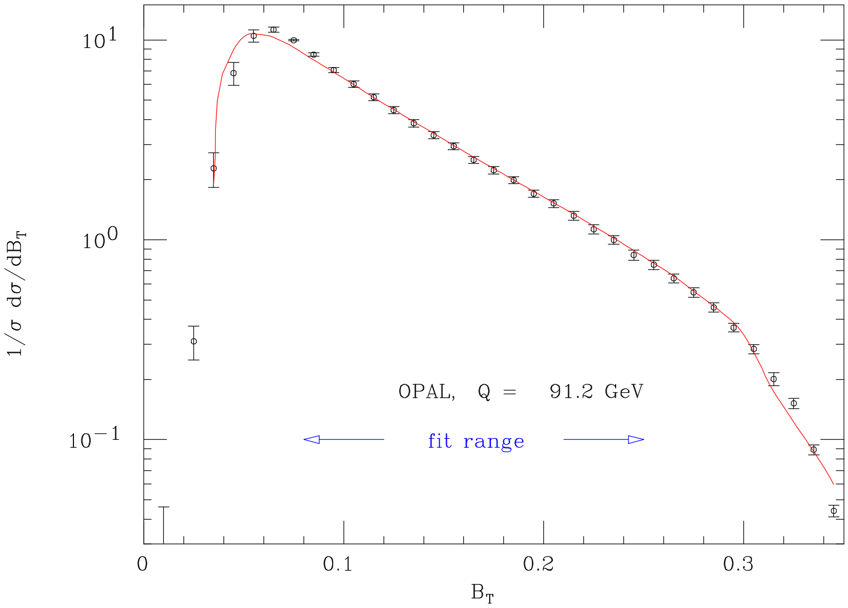
$$B = \frac{\sum_i |\boldsymbol{p}_i \times \boldsymbol{n}_T|}{2\sum_i |\boldsymbol{p}_i|}$$

$$\frac{1}{\sigma}\frac{d\sigma}{dB} \simeq f_B^{\text{pert}}\left(B - \frac{D_B}{Q}\ln\frac{B_0}{B}\right)$$

where $D_B \simeq 0.5$ GeV.

YuL Dokshitzer, G Marchesini & GP Salam, hep-ph/9812487

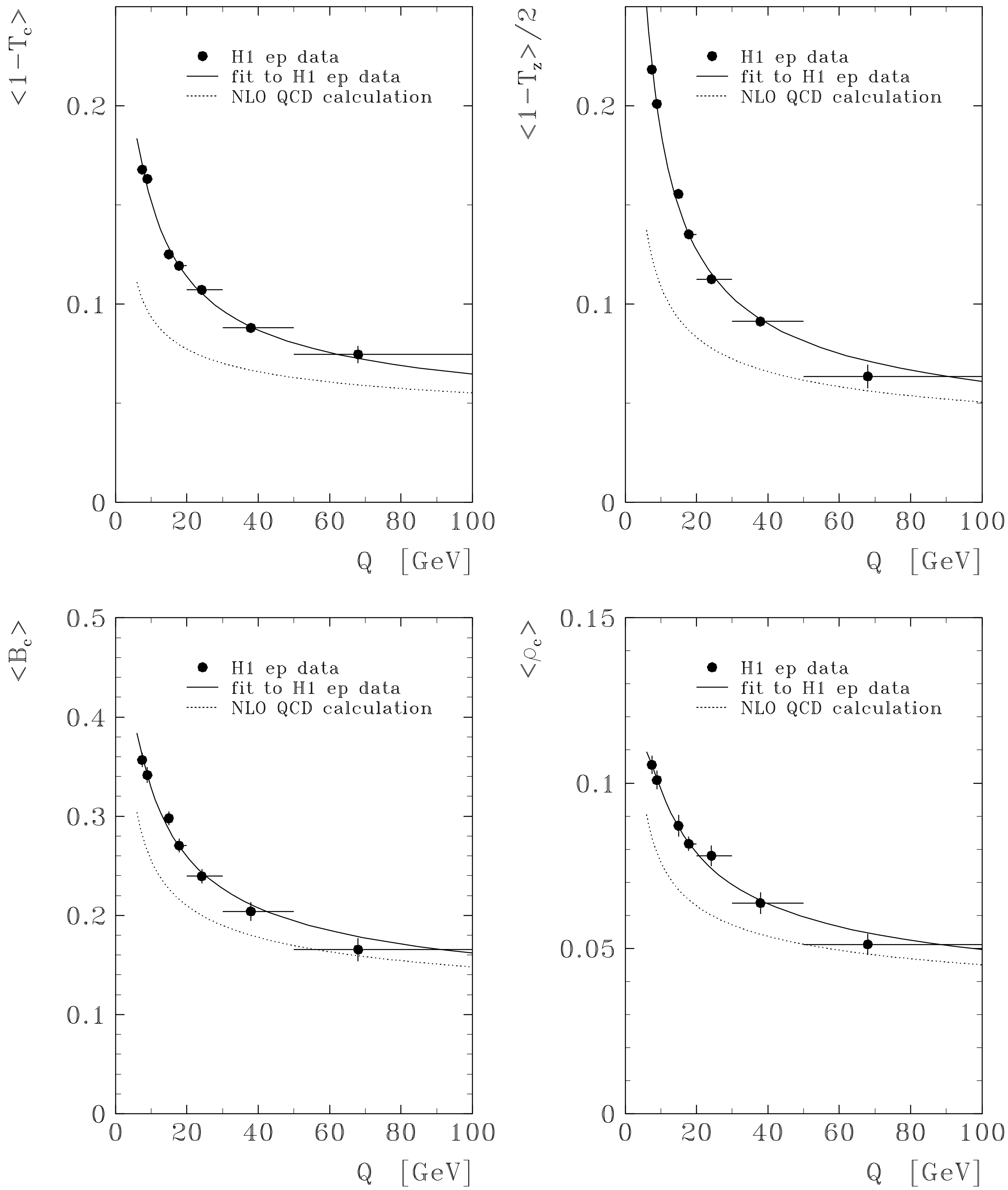




DIS event shapes

Similar results (and similar values of α_0) are obtained from event shapes in deep inelastic lepton scattering:

H1 Collaboration, C Glasman et al., hep-ex/9706002



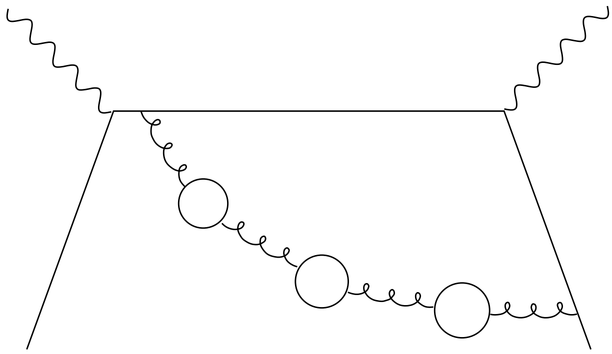
Asymptotics of perturbation theory

It is expected that perturbation series will diverge at high order.

There may even be signs of this already for some observables, e.g. the Gross–Llewellyn Smith sum rule:

$$\frac{1}{6}\int_0^1 dx\,(F_3^{\nu}+F_3^{\bar\nu}) = 1-\frac{\alpha_S}{\pi}-3.6\left(\frac{\alpha_S}{\pi}\right)^2-19\left(\frac{\alpha_S}{\pi}\right)^3+\cdots$$

One source of divergence is *renormalon chain* graphs



$$F \sim \alpha_S \sum_{n=0}^{\infty} n!\left(\frac{\beta_0\alpha_S}{2\pi p}\right)^n$$

- Quark loops $\Rightarrow \beta_0 = -2N_f/3$

  $p = 1, 2, \ldots$ depends on observable F

  Infrared renormalon $\Leftrightarrow$ low momentum in loops

- "Naive non-Abelianization": replace

$$N_f \to N_f - 33/2\,, \qquad \beta_0 \to 11 - 2N_f/3$$

Includes gluon loops plus . . .

Seems to be good even for $N_f = 0$ (lattice)

- If perturbation series is *asymptotic* then

$$F \sim \alpha_S \sum n! \left(\frac{\beta_0 \alpha_S}{2\pi p}\right)^n \equiv \sum f_n$$

implies

$$F = \sum^N f_n \pm f_N$$

Truncating at smallest term, $f_N = \min\{f_n\}$, gives

$$N \sim \frac{2\pi p}{\beta_0 \alpha_S}\,, \qquad \delta F = f_N \sim \exp\left(-\frac{2\pi p}{\beta_0 \alpha_S}\right)$$

$$\alpha_S(Q^2) \sim \frac{4\pi}{\beta_0 \ln(Q^2/\Lambda^2)} \qquad \Rightarrow \qquad \delta F \sim \left(\frac{\Lambda}{Q}\right)^p$$

- This *renormalon ambiguity* must be cancelled by a non-perturbative contribution with the same power behaviour

‘Dispersive’ approach

- Assume that $\alpha_S(\mu^2)$ can be defined more generally (and universally?).

  Then it should have singularities for real –ve μ^2 only, and we can write a *dispersion relation*

$$\alpha_S(q^2) = -\int_0^\infty \frac{d\mu^2}{\mu^2+q^2}\rho_S(\mu^2)$$

  where ρ_S is the *discontinuity* of α_S:

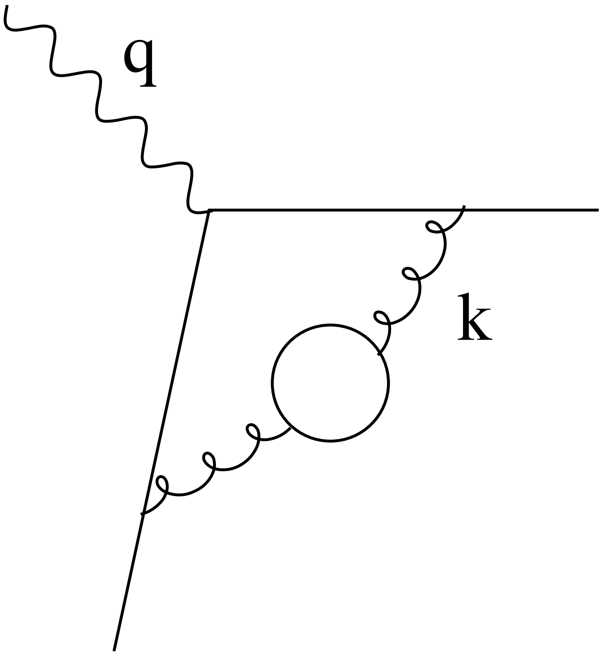
$$\begin{aligned}\rho_S(\mu^2) &= \frac{1}{2\pi i}\left[\alpha_S\left(\mu^2 e^{i\pi}\right) - \alpha_S\left(\mu^2 e^{-i\pi}\right)\right] \\ &\equiv \frac{1}{2\pi i}\mathrm{Disc}[\alpha_S(-\mu^2)] \qquad \left(\sim \mu^2\frac{d\alpha_S}{d\mu^2}\right)\end{aligned}$$

- Replacing a gluon by a renormalon chain gives

$$\begin{aligned}F(Q^2) &= \alpha_S(0)\mathcal{F}(0) + \int_0^\infty \frac{d\mu^2}{\mu^2}\rho_S(\mu^2)\mathcal{F}\left(\frac{\mu^2}{Q^2}\right) \\ &= \int_0^\infty \frac{d\mu^2}{\mu^2}\rho_S(\mu^2)\left[\mathcal{F}\left(\frac{\mu^2}{Q^2}\right) - \mathcal{F}(0)\right]\end{aligned}$$

  where $\mathcal{F}(\mu^2/Q^2)$ is the first-order expansion coefficient for the observable $F(Q^2)$ *with the gluon mass set equal to* μ.

- Consider for example the vertex correction



$$
\begin{aligned}
F_{\mathsf{V}}(Q^2) &= \int_{-\infty}^{0} \frac{dk^2}{k^2} f\left(\frac{k^2}{Q^2}\right) \alpha_{\mathsf{S}}(-k^2) \\
&= \int_{-\infty}^{0} \frac{dk^2}{k^2} f\left(\frac{k^2}{Q^2}\right) \int_{0}^{\infty} \frac{d\mu^2}{k^2-\mu^2} \rho_{\mathsf{S}}(\mu^2) \\
&= \int_{0}^{\infty} \frac{d\mu^2}{\mu^2} \rho_{\mathsf{S}}(\mu^2) \int_{-\infty}^{0} \left[\frac{dk^2}{k^2-\mu^2} - \frac{dk^2}{k^2}\right] f\left(\frac{k^2}{Q^2}\right) \\
&\equiv \int_{0}^{\infty} \frac{d\mu^2}{\mu^2} \rho_{\mathsf{S}}(\mu^2) \left[\mathcal{F}_{\mathsf{V}}\left(\frac{\mu^2}{Q^2}\right) - \mathcal{F}_{\mathsf{V}}(0)\right]
\end{aligned}
$$

Perturbative contribution: scale setting

$$F(Q^2) = \int_0^\infty \frac{d\mu^2}{\mu^2}\rho_{\mathrm{S}}(\mu^2)\left[\mathcal{F}\left(\frac{\mu^2}{Q^2}\right) - \mathcal{F}(0)\right]$$

- How does this relate to the standard perturbative result? Using $\rho_{\mathrm{S}}(\mu^2) \sim \mu^2\, d\alpha_{\mathrm{S}}/d\mu^2$ we have

$$\begin{aligned} F(Q^2) &\sim \int_0^\infty d\mu^2 \frac{d\alpha_{\mathrm{S}}}{d\mu^2}\left[\mathcal{F}\left(\frac{\mu^2}{Q^2}\right) - \mathcal{F}(0)\right] \\ &= -\int_0^\infty d\epsilon\, \alpha_{\mathrm{S}}(\epsilon Q^2)\,\frac{d\mathcal{F}}{d\epsilon} \end{aligned}$$

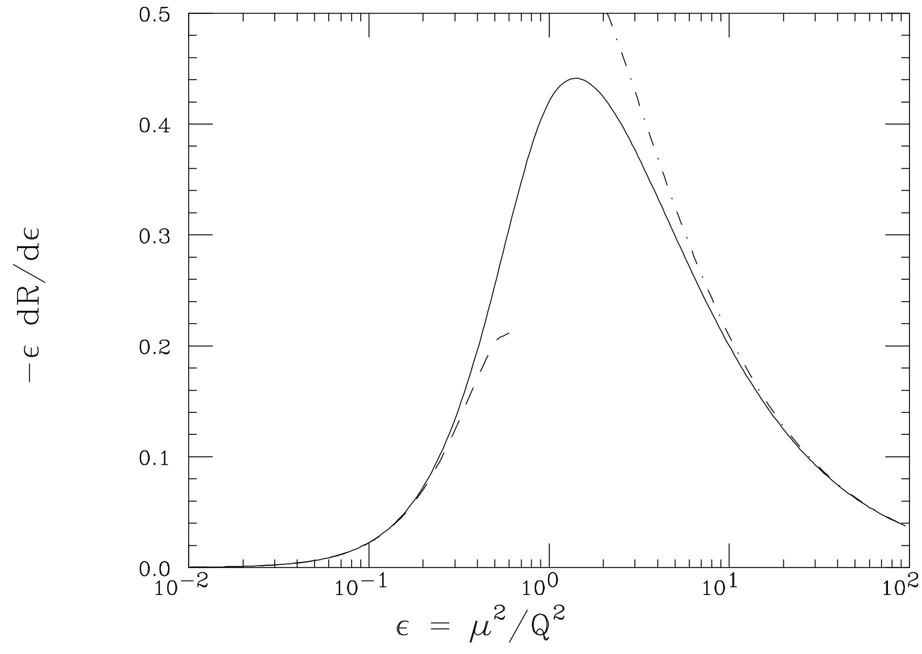
If $d\mathcal{F}/d\epsilon$ peaks around $\epsilon = \xi$ we can write

$$F(Q^2) \simeq -\alpha_{\mathrm{S}}(\xi Q^2)\int_0^\infty d\epsilon\,\frac{d\mathcal{F}}{d\epsilon} = \alpha_{\mathrm{S}}(\xi Q^2)\,\mathcal{F}(0)$$

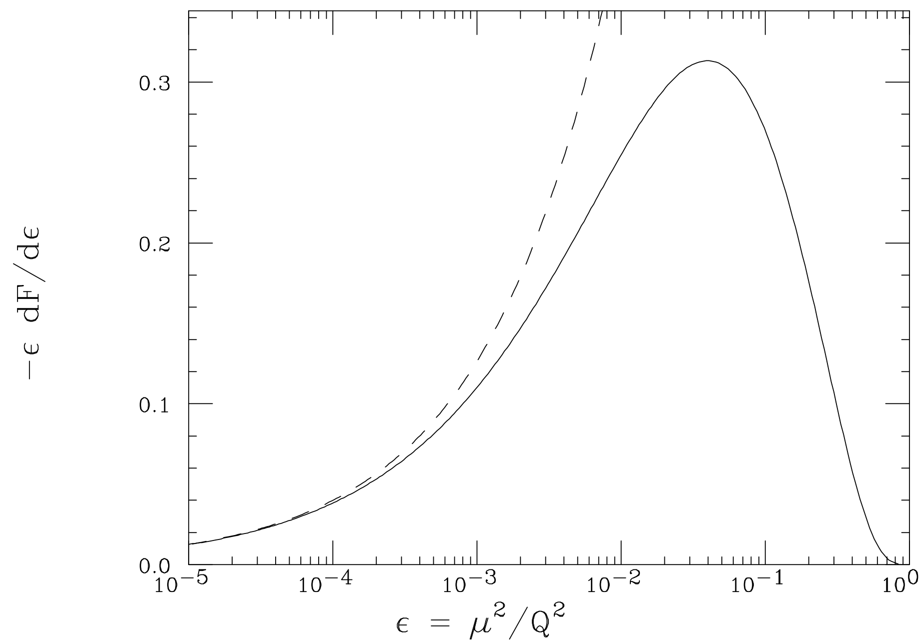
- Thus the *peak position* in $d\mathcal{F}/d\epsilon$ determines the optimal choice of scale for α_{S}.

- Peak is at $\epsilon = \xi \simeq 1$ for $F = R_{e^+e^-}$, the e^+e^- total cross section:

YuL Dokshitzer, G Marchesini & BRW, hep-ph/9512336



- However, $\xi \ll 1$ for event shapes, e.g. $\langle 1 - T \rangle$:



This explains the large higher-order corrections when such quantities are expressed in terms of $\alpha_S(Q^2)$.

Non-perturbative contribution

Now suppose that $\alpha_{\mathsf{S}}(\mu^2)$ has some non-perturbative modification $\delta\alpha_{\mathsf{S}}(\mu^2)$ at low μ^2. The corresponding non-perturbative contribution to $F(Q^2)$ will be

$$\delta F = \frac{1}{2\pi i}\int_0^\infty \frac{d\mu^2}{\mu^2}\mathrm{Disc}[\delta\alpha_{\mathsf{S}}(-\mu^2)]\left[\mathcal{F}\left(\frac{\mu^2}{Q^2}\right) - \mathcal{F}(0)\right]$$

- Rotating the contour by $\mp\pi$ in the two terms of the discontinuity gives

$$\begin{aligned}\delta F &= \int_0^\infty \frac{d\mu^2}{\mu^2}\delta\alpha_{\mathsf{S}}(\mu^2)\,\mathcal{G}\left(\frac{\mu^2}{Q^2}\right)\\ \mathcal{G}(\epsilon) &= -\frac{1}{2\pi i}\mathrm{Disc}[\mathcal{F}(-\epsilon)]\end{aligned}$$

- Since $\delta\alpha_{\mathsf{S}}(\mu^2)$ is confined to small μ^2 (say $\mu^2 < \mu_{\mathsf{I}}^2$), the leading Q^2-dependence of δF is determined by the $\epsilon \longrightarrow 0$ behaviour of $\mathcal{F}(\epsilon)$.

  In particular, only non-analytic behaviour at $\epsilon = 0$ can contribute to δF

- For event shapes $\mathcal{F}(\epsilon) \sim \mathcal{F}(0) - a_F\sqrt{\epsilon}$, which implies

$$\delta F = \frac{2a_F}{\pi}\frac{1}{Q}\int d\mu\,\delta\alpha_{\mathsf{S}}(\mu^2)$$

Universal strong coupling

Since we give real significance to $\alpha_{\mathsf{S}}(\mu^2)$ at low μ^2, we have to define a preferred renormalization scheme.

- We use the *bremsstrahlung scheme*, in which the divergent part of the soft gluon spectrum is absorbed into the running coupling:

$$\alpha_{\mathsf{B}} \sim \alpha_{\overline{\mathsf{MS}}} + \frac{K_{\mathsf{B}}}{2\pi}\alpha^2_{\overline{\mathsf{MS}}}$$

where $K_{\mathsf{B}} = 67/6 - \pi^2/2 - 5N_f/9$. This is close to the *heavy quark potential* or V-scheme, which has $K_{\mathsf{V}} = 31/6 - 5N_f/9$

- If F^{pert} is computed to second order in $\alpha_{\overline{\mathsf{MS}}} \equiv \alpha_{\overline{\mathsf{MS}}}(Q^2)$ we therefore have

$$\begin{aligned}
\delta\alpha_{\mathsf{S}}(\mu^2) &\simeq \alpha_{\mathsf{B}}(\mu^2) - \left[\alpha_{\overline{\mathsf{MS}}} + \frac{\beta_0}{2\pi}\left(\ln\frac{Q}{\mu} + \frac{K_{\mathsf{B}}}{\beta_0}\right)\alpha^2_{\overline{\mathsf{MS}}}\right] \\
&\simeq 0 \qquad \text{for } \mu > \mu_{\mathsf{I}}
\end{aligned}$$

Thus

$$\delta F = -\frac{2a_F}{\pi}\frac{\mu_{\mathsf{I}}}{Q}\left[\alpha_0(\mu_{\mathsf{I}}) - \alpha_{\overline{\mathsf{MS}}} - \frac{\beta_0}{2\pi}\left(\ln\frac{Q}{\mu_{\mathsf{I}}} + \frac{K_{\mathsf{B}}}{\beta_0} + 1\right)\alpha^2_{\overline{\mathsf{MS}}}\right]$$

where

$$\alpha_0(\mu_{\mathsf{I}}) \equiv \frac{1}{\mu_{\mathsf{I}}}\int_0^{\mu_{\mathsf{I}}} d\mu\, \alpha_{\mathsf{B}}(\mu^2)$$

- From experiment,

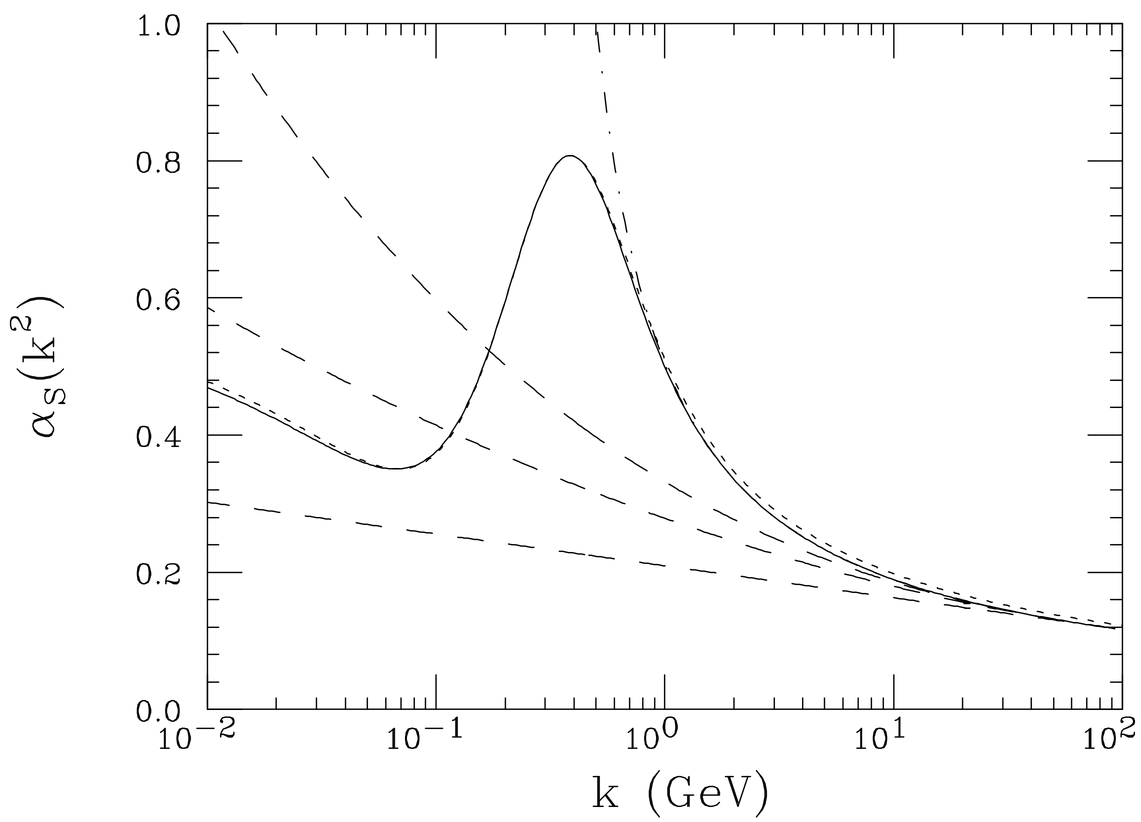
$$\alpha_0(\mu_I) \simeq 0.5 \qquad \text{for } \mu_I = 2 \text{ GeV}$$

Simple models for α_B can be constructed to give this value, for example

$$\alpha_B(\mu^2) = \frac{4\pi}{\beta_0} a\left(\frac{\mu^2}{\Lambda^2}\right)$$

$$a(z) = \frac{1}{\ln z} + 125\frac{1+4z}{(1-z)(4+z)^4}$$

BRW, hep-ph/9805484



Results on event shapes

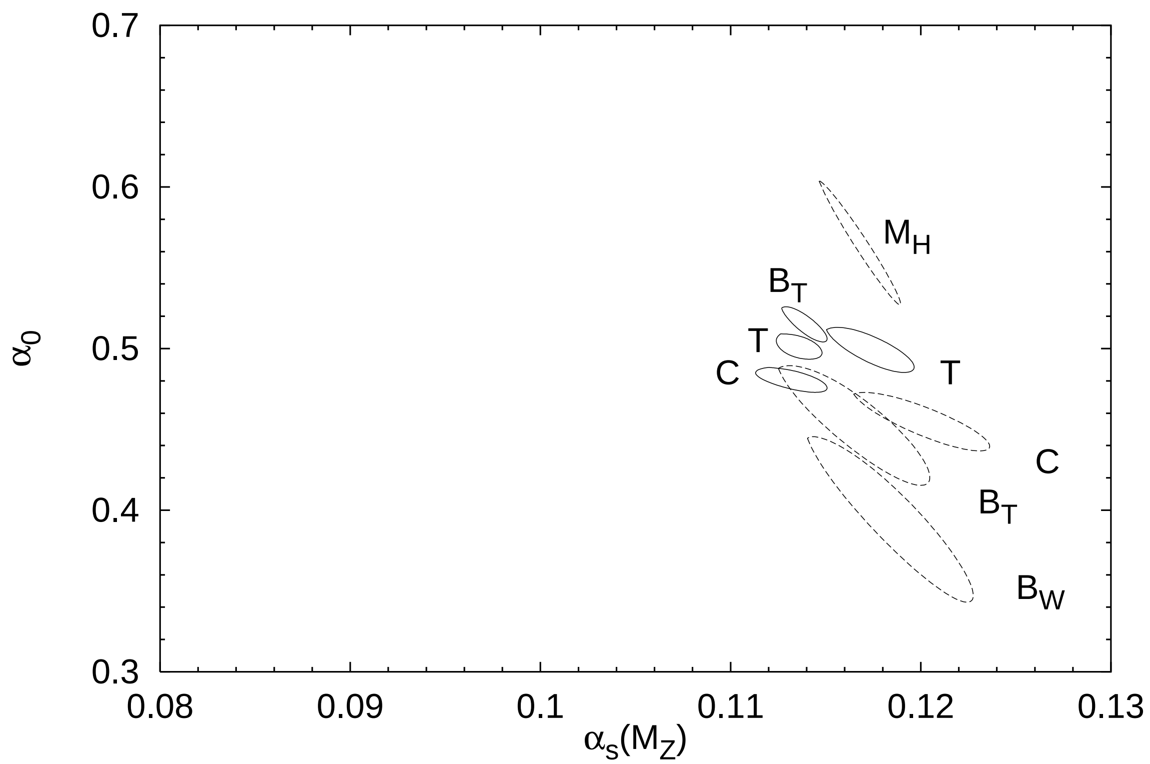
The coefficients a_F for the $1/Q$ corrections to event shapes have been computed to two-loop order:

F	$1-T$	C	M_T^2/Q^2	M_H^2/Q^2	B_T	B_W
a_F	3.1	14.4	3.1	1.5	1.5	0.8

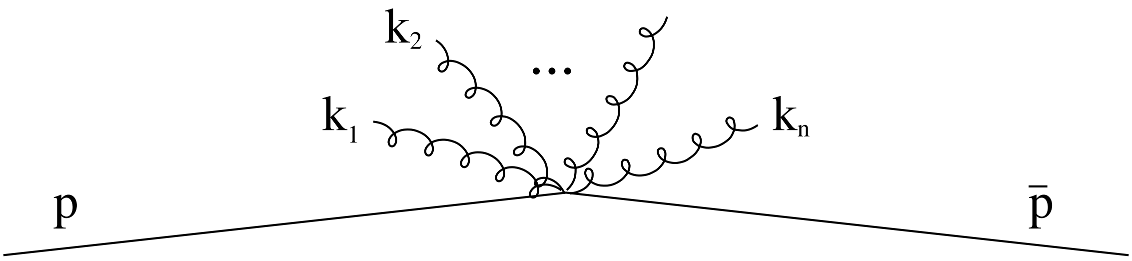
Using these one can make fits to $\alpha_{\overline{\rm MS}}$ and α_0:

Variable	$\alpha_{\overline{\rm MS}}$	α_0	χ^2/d.o.f.
$1-T$	0.1177 ± 0.0013	0.498 ± 0.009	$57.0/38$
C	0.1206 ± 0.0021	0.453 ± 0.011	$10.7/8$
M_H^2/Q^2	0.1171 ± 0.0012	0.560 ± 0.022	$15.2/25$
B_T	0.1170 ± 0.0023	0.451 ± 0.023	$14.9/21$
B_W	0.1189 ± 0.0025	0.391 ± 0.031	$12.8/20$

YuL Dokshitzer, G Marchesini & GP Salam, hep-ph/9812487



- Outline (1–loop) calculation of a_F for $t \equiv 1 - T$:



Sudakov decomposition of gluon momenta k_i:

$$k_i = \alpha_i p + \beta_i \bar{p} + k_{ti}$$

where $p \cdot k_{ti} = \bar{p} \cdot k_{ti} = 0$, $k_{ti}^2 \equiv -\boldsymbol{k}_i^2$.

In the soft region quark/antiquark recoils are negligible,

$$t \simeq \sum_i \min\{\alpha_i, \beta_i\}$$

and the cross section approximately factorizes:

$$\frac{d\sigma_n}{\sigma_0} \simeq \prod_{i=1}^{n} \frac{C_F}{\pi^2} \frac{d\alpha_i}{\alpha_i} \frac{d\beta_i}{\beta_i} d^2\boldsymbol{k}_i \, \alpha_{\rm S}(\boldsymbol{k}_i^2) \delta(\alpha_i \beta_i Q^2 - \boldsymbol{k}_i^2)$$

Defining

$$\Sigma(t) = \int_0^t dt \, \frac{1}{\sigma} \frac{d\sigma}{dt}$$

and using

$$\Theta(t) = \int_{\frac{1}{2}-i\infty}^{\frac{1}{2}+i\infty} \frac{d\nu}{2i\pi\nu} e^{\nu t}$$

we can write

$$\Sigma(t) \simeq \int_{\frac{1}{2}-i\infty}^{\frac{1}{2}+i\infty} \frac{d\nu}{2i\pi\nu} e^{\nu t - R(\nu)}$$

where

$$R(\nu) = \frac{C_F}{\pi} \int \frac{d\alpha}{\alpha}\frac{d\beta}{\beta} d\boldsymbol{k}^2 \alpha_{\mathsf{S}}(\boldsymbol{k}^2)[1 - e^{-\nu \min\{\alpha,\beta\}}]\delta(\alpha\beta Q^2 - \boldsymbol{k}^2)$$

- A non-perturbative correction $\delta\alpha_{\mathsf{S}}(\boldsymbol{k}^2)$ at small $\boldsymbol{k}^2$ therefore generates a *shift* in the distribution of t:

$$\Sigma(t) \simeq \Sigma^{\mathsf{pert}}(t - \delta t)$$

where

$$\begin{aligned} \delta t = \frac{\delta R}{\nu} &= \frac{C_F}{\pi} \int \frac{d\alpha}{\alpha}\frac{d\beta}{\beta} d\boldsymbol{k}^2\, \delta\alpha_{\mathsf{S}}(\boldsymbol{k}^2) \min\{\alpha,\beta\} \\ &\quad \cdot \delta(\alpha\beta Q^2 - \boldsymbol{k}^2) \\ &= \frac{4C_F}{\pi}\frac{1}{Q} \int d|\boldsymbol{k}|\, \delta\alpha_{\mathsf{S}}(\boldsymbol{k}^2) \end{aligned}$$

Hence in this approximation $a_T = 2C_F = 8/3$.

Results on structure functions

- Non-flavour-singlet

The non-singlet quark coefficient functions have $\epsilon \to 0$ behaviour

$$\mathcal{F}(x,\epsilon) \sim -P(x)\ln\epsilon + C(x) + D(x)\epsilon\ln\epsilon$$

This has to be convoluted with the quark distribution $q(x,Q^2)$. The first term gives the usual logarithmic scaling violation. The leading power correction is given by the $\epsilon\ln\epsilon$ term:

$$\delta F(x) = \frac{1}{Q^2}\int d\mu^2\,\delta\alpha_{\mathrm{S}}(\mu^2)\,D(x)\otimes q(x,Q^2)$$

The curves shown earlier correspond to

$$\int d\mu^2\,\delta\alpha_{\mathrm{S}}(\mu^2) \simeq 1\ \mathrm{GeV}^2$$

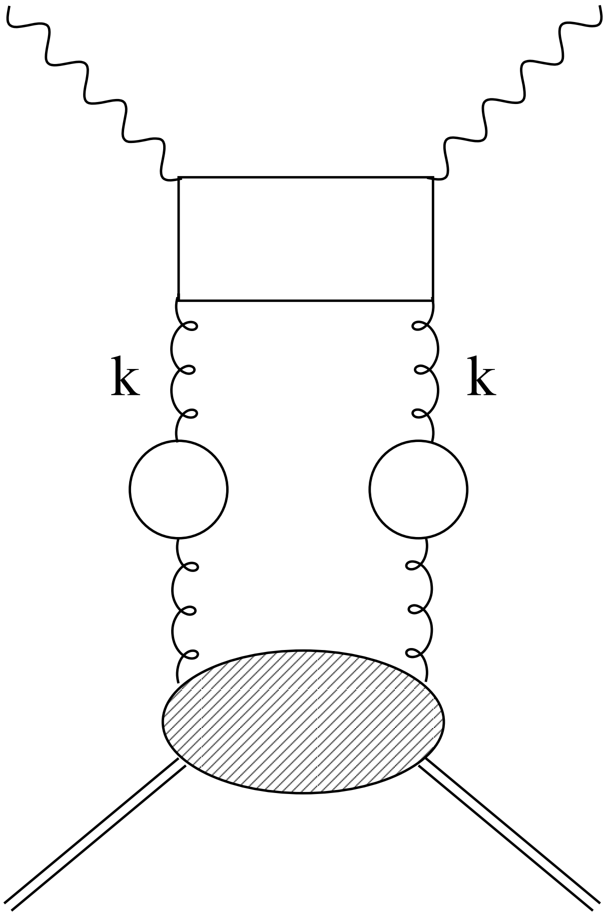
The dominant term in $D(x)$ at large x is $C_F\delta'(1-x)/2\pi$, which gives

$$D(x)\otimes q(x,Q^2) \simeq -\frac{C_F}{2\pi}\frac{\partial}{\partial x}x\,q(x,Q^2)$$

Hence the $1/Q^2$ correction rises sharply at high x.

Flavour singlet

Flavour singlet structure functions have leading contributions involving two renormalon chains:



- The dispersive method now gives

$$
\begin{aligned}
F \quad = \quad & \int \frac{d\mu_1^2}{\mu_1^2}\frac{d\mu_2^2}{\mu_2^2}\rho_{\rm S}(\mu_1^2)\rho_{\rm S}(\mu_2^2)\left[\mathcal{F}\left(\frac{\mu_1^2}{Q^2},\frac{\mu_2^2}{Q^2}\right)\right. \\
& \left. -\mathcal{F}\left(\frac{\mu_1^2}{Q^2},0\right)-\mathcal{F}\left(0,\frac{\mu_2^2}{Q^2}\right)+\mathcal{F}(0,0)\right]
\end{aligned}
$$

❖ Since both chains have the *same momentum* k, $\mathcal{F}$ has the form

$$\begin{aligned}\mathcal{F}(\epsilon_1, \epsilon_2) &= \int_0^\infty d\kappa \frac{f(\kappa)}{(\kappa+\epsilon_1)(\kappa+\epsilon_2)} \\ &= \frac{\epsilon_1\mathcal{F}(\epsilon_1) - \epsilon_2\mathcal{F}(\epsilon_2)}{\epsilon_1 - \epsilon_2}\end{aligned}$$

where $\epsilon_i = \mu_i^2/Q^2$, $\kappa = -k^2/Q^2$ and

$$\mathcal{F}(\epsilon) = -\frac{1}{\epsilon}\int_0^\infty d\kappa \frac{f(\kappa)}{\kappa+\epsilon}$$

❖ After some manipulation we find that the non-perturbative contribution is

$$\delta F = \int_0^\infty \frac{d\mu^2}{\mu^2}\delta[\alpha_{\mathsf{S}}(\mu^2)]^2\, \mathcal{G}\left(\frac{\mu^2}{Q^2}\right)$$

where

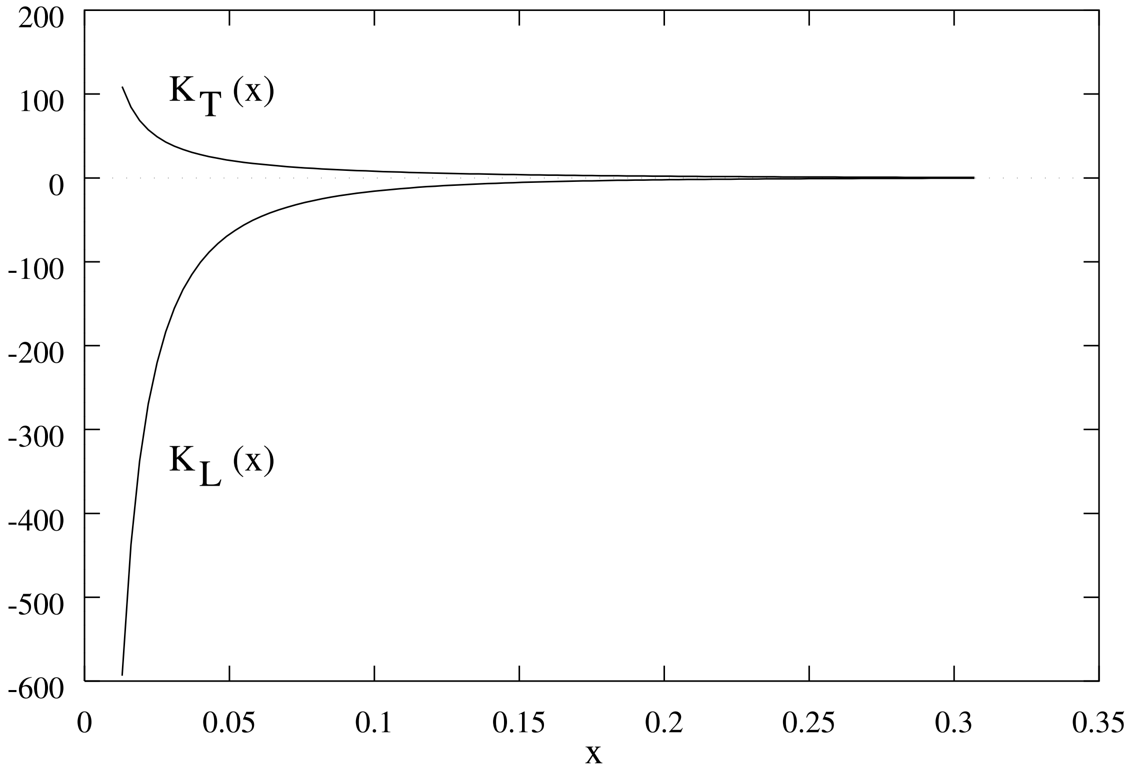
$$\mathcal{G}(\epsilon) = -\frac{1}{2\pi i}\mathsf{Disc}[\mathcal{F}(-\epsilon)]$$

$$\delta[\alpha_{\mathsf{S}}(\mu^2)]^2 \simeq [\alpha_{\mathsf{B}}(\mu^2)]^2 - \left[\alpha_{\overline{\mathsf{MS}}} + \frac{\beta_0}{2\pi}\left(\ln\frac{Q}{\mu} + \frac{K_{\mathsf{B}}}{\beta_0}\right)\alpha^2_{\overline{\mathsf{MS}}}\right]^2$$

❖ The resulting power corrections have the form

$$\delta F_i(x) = \frac{C_F T_R}{(2\pi)^2} \frac{K_i(x)}{Q^2} \int d\mu^2 \, \delta[\alpha_S(\mu^2)]^2$$

GE Smye, hep-ph/9810292



Conclusions

- Power corrections are present, and must be taken into account when comparing perturbative QCD with data
- The observed powers are consistent with those deduced from renormalon/dispersive analyses
- At least such analyses give an indication of which observables are best for measuring $\alpha_S(M_Z^2)$
- The hypothesis of a universal strong coupling at low scales seems to work surprisingly well
- The mean value of the coupling deduced from $1/Q$ corrections is not large: $\alpha_0(2\ \text{GeV}) \simeq 0.5$
- Still many outstanding questions:
  - More rigorous formulation (pinch technique?)
  - Connection with confinement/hadronization?
  - Connection with quark/gluon condensates?
  - Power corrections to α_S itself?
  - Shape variables without $1/Q$ corrections?
  - Treatment of multi-chains, higher orders
  - Tests for singlet quantities