Hadron interactions, color and QCD partons

3. Probing non-perturbative phenomena with perturbative tools

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Screwing Non-Perturbative QCD with Perturbative Tools

About 15 years ago first theoretical attempts have been made to quantify genuine non-perturbative effects in perturbatively calculable (CIS) observables

The test case : the total cross section of e^+e^- annihilation into hadrons.

To predict $\sigma_{tot} \rightarrow$ hadrons one calculates instead the cross sections of quark and gluon production, $(e^+ e^- \rightarrow q\,\bar{q}) + (e^+ e^- \rightarrow q\,\bar{q} + g)$ + etc., where quarks and gluons are being treated ***perturbatively*** as real (un-confined, flying) objects.

The ***completeness*** argument provides an apology for such a brave substitution :

> Once instantaneously produced by the electromagnetic (electroweak) current, the quarks (and secondary gluons) have nowhere else to go but to convert, ***with unit probability***, into hadrons in the end of the day.

This ***guess*** looks rather solid and sounds convincing, but relies on two hidden assumptions:

1. The allowed hadron states should be ***numerous*** as to provide the quark-gluon system the means for "**regrouping**", "**blanching**", "**fitting**" into hadrons.

2. It implies that the "**production**" and "**hadronization**" stages of the process can be separated and treated independently.

1. To comply with the first assumption ("*regrouping*") the annihilation energy has to be taken large enough, $s \equiv Q^2 \gg s_0$. In particular, it fails miserably in the resonance region $Q^2 < s_0 \sim 2M^2_{res}$.

Thus, the point-by-point correspondence between hadron and quark cross sections, $\sigma^{tot}_{hadr}(Q^2)$? = $\sigma^{tot}_{qq}(Q^2)$, cannot be sustained except at very high energies.

It can be traded, however, for something more manageable.

Invoking the dispersion relation for the photon propagator (causality = analyticity) one can relate the ***energy integrals*** of $\sigma^{tot}(s)$ with the correlator of electromagnetic currents in a deeply Euclidean region of large ***negative*** Q^2.

The latter corresponds to small space-like distances between interaction points, where the perturbative approach is definitely valid.

Expanding the answer in a formal series of local operators, one arrives at the structure in which the corrections to the trivial unit operator generate the usual perturbative series in powers of α_s (logarithmic corrections), whereas the vacuum expectation values of dimension-full (Lorentz- and colour-invariant) QCD operators provide non-perturbative corrections suppressed as powers of Q.

This is the realm of the famous "**ITEP sum rules**" which proved to be successful in linking the parameters of the low-lying resonances in the ***Minkowsky space*** with expectation values characterising a non-trivial structure of the QCD vacuum in the ***Euclidean space***.

ITEF "NP physics" = ***non-singular long range gluon fields***

Shifman, Vainstein & Zakharov *Nucl Phys B*. (1979)

The leaders among them are the gluon condensate **$\langle \alpha_s G_{\mu\nu} G^{\mu\nu} \rangle$** and the quark condensate $\langle\psi\bar{\psi}\rangle\langle\psi\bar{\psi}\rangle$ which contribute to the total annihilation cross section, symbolically, as

$$\sigma^{\text{tot}}_{\text{hadr}}(Q^2) - \sigma^{\text{tot}}_{q\bar{q}+X}(Q^2) = c_1 \frac{\alpha_s G^2}{Q^4} + c_2 \frac{\langle\psi\bar{\psi}\rangle^2}{Q^6} + \dots$$

2. Validating the second assumption ("*production vs. hadronization*") also calls for large **Q^2**. To be able to separate the two stages of the process, it is *necessary* to have the production time of the quark pair **Q^{-1}** to be much smaller than the time **$t_1 \sim \mu^{-1} \sim 1$ fm/c** when the first hadron appears in the system. Whether this condition is *sufficient,* is another valid question. And a tricky one.

As we know, due to the gluon bremsstrahlung the perturbative production of secondary gluons and quark pairs spans an immense interval of time, ranging from a very short time, **$t_{\text{form}} \sim Q^{-1} \ll t_1$**, all the way up to a macroscopically large time **$t_{\text{form}} \sim Q/\mu^2 \gg t_1$**.

This accompanying radiation is responsible for formation of ***hadron jets***.

It does not, however, affect the ***total cross section***. It is the rare hard gluons with large energies and transverse momenta, **$\sim Q$**, that only matter.

This follows from the celebrated ***Bloch-Nordsieck theorem*** which states that the logarithmically enhanced (divergent) contributions due to real production of *collinear* and *soft* quanta cancel against the corresponding virtual corrections :

$$\sigma^{\text{tot}}_{q\bar{q}+X} = \sigma_{Born}\left(1 + \frac{\alpha_s}{\pi}\left[\infty_{\text{real}} - \infty_{\text{virtual}}\right] + \dots\right) = \sigma_{Born}\left(1 + \frac{3}{4}\frac{C_F \alpha_s(Q^2)}{\pi} + \dots\right)$$

"massive gluon"

Can the Bloch-Nordsieck result hold ***beyond*** perturbation theory?

Looking into this problem produced an extremely interesting result that has laid a foundation for the development of perturbative techniques aimed at analysing non-perturbative effects.

V. Braun, M. Beneke and V. Zakharov have demonstrated that the real-virtual cancellation actually proceeds ***much deeper*** than was originally expected.

[*Phys.Rev.Lett.* ***73*** *(1994) 3058*]

- Introduce into the calculation of the radiative correction ***gluon mass*** m as an IR cutoff.

- Study the dependence of the answer on m.

  A CIS quantity, by definition, remains finite in the limit $m=0$. This does not mean, however, that it is totally **insensitive** to the modification of gluon propagation.

  In fact, the m-dependence provides a handle for analysing the *small transverse momenta* inside Feynman integrals. It is this region of integration over parton momenta where the QCD coupling gets out of control and the genuine NP physics comes onto the stage.

- Then, the sensitivity of a given CIS observable to the ***infrared domain*** is determined by the first non-vanishing term ***non-analytic*** in m^2 at $m=0$.

Bloch-Nordsieck theorem extended

In the case of one-loop analysis of the total annihilation cross section that we are discussing, one finds that in the sum of real and virtual contributions not only the terms singular at ***m=0***,

$$\ln^2 m^2 \quad \text{and} \quad \ln m^2 \,,$$

cancel, *as required by the Bloch-Nordsieck theorem*,

but that the cancellation extends also to the whole tower of ***finite terms***

$$m^2 \ln^2 m^2,\ m^2 \ln m^2,\ m^2,\ m^4 \ln^2 m^2,\ m^4 \ln m^2 \,.$$

The first ***non-analytic term*** appears at the level of m^6 :

$$\frac{3}{4}\frac{C_F\alpha_s}{\pi}\left(1+2\frac{m^6}{Q^6}\ln\frac{m^2}{Q^2}+\mathcal{O}\left(m^8\right)\right)$$

It signals the presence of the non-perturbative Q^{-6} correction, which is equivalent to that of the ITEP ***quark condensate***.
The ***gluon condensate*** contribution emerges in the next order in α_s

Why ***"gluon mass"*** ? What is it and how does it serve as a large-distance probe ?

wave function renormalization $\Longrightarrow$ ***running coupling***

$$G(k^2)=\frac{1}{-k^2-i0} \Longrightarrow \frac{Z(k^2)}{-k^2-i0} = \frac{\alpha_s(-k^2)}{-k^2-i0} \qquad \frac{Z(-k^2)}{Z(-\mu_R^2)}=\frac{\alpha(k^2)}{\alpha(\mu_R^2)}$$

We want this identification to make sense in the entire k^2 plane

We know sufficiently well how α_s**(-k²)** behaves in the Euclidean region, at large negative **k²** while we know next to nothing about the small-**k²** region.

However, whatever the function α_s is, it had better respect ***causality***.

Therefore we suppress the formal PT *tachion* ("Landau singularity") and choose the "***physical cut***" alone, **0<k²**, as a support for the **dispersive relation** :

$$\boxed{\alpha_s(q^2) = \int_0^\infty \frac{dm^2\, q^2}{(m^2+q^2)^2}\,\alpha_{\rm eff}(m^2)}$$

$$\frac{d}{d\ln\mu^2}\alpha_{\rm eff}(\mu^2) = -\frac{1}{2\pi i}\,\mathrm{Disc}\,\{\alpha_s(-\mu^2)\}$$

It can be formally inverted as an operator relation :

$$\alpha_{\rm eff}(\mu^2) = \frac{\sin(\pi\mathcal{P})}{\pi\mathcal{P}}\,\alpha_s(\mu^2)\,, \quad \mathcal{P} = \mu^2\frac{d}{d\mu^2}$$

We are ready now for a "***heavy gluon***" :

$$\boxed{\frac{\alpha_s(-k^2)}{-k^2-i0} = \int_0^\infty \frac{dm^2}{m^2}\,\alpha_{\rm eff}(m^2)\cdot\frac{-d}{d\ln m^2}\frac{1}{m^2-k^2-i0}}$$

Substitute into the Feynman diagram, and integrate first over the gluon 4-momentum :

$$\boxed{V(Q^2,x) = \int_0^\infty \frac{dm^2}{m^2}\alpha_{\rm eff}(m^2)\dot{\mathcal{F}}_V(\epsilon,x)}$$

$$\epsilon = m^2/Q^2 \qquad \dot{\mathcal{F}} \equiv \frac{-d\mathcal{F}}{d\ln\epsilon}$$

$\boldsymbol{\mathcal{F}_V}$ - a "***Characteristic Function***" for the observable ***V***

perturbative answer

At this point there is no difference with the usual PT answer.

The CIS nature of the observable V guarantees convergence of the m^2 integration :

$\dot{\mathcal{F}}$ vanishes as a power of $\epsilon\,(\epsilon^{-1})$ in the $\epsilon \to 0\ (\epsilon \to \infty)$ limit.

Therefore the distribution $\dot{\mathcal{F}}$ ***has a maximum*** at some $\epsilon = C_V(x) = \mathcal{O}(1)$,
and the integral is dominated by the large-momentum region $m^2 \sim Q^2$.

Approximating $\alpha_{\rm eff}(\mathbf{m^2}) \simeq \alpha_{\mathbf{s}}(\mathbf{Q^2})$ we reproduce the one-loop PT answer :

$$V(Q^2,x) \simeq \alpha_s(Q^2)\mathcal{F}_V(0,x)$$

Using the *observable-dependent position of the maximum* of the m^2-distribution as the scale for the coupling, $\alpha_{\mathbf{s}}(\mathbf{C_V(x)\cdot Q^2})$, does a ***better job*** since it minimizes higher order effects.

The dispersive technology in this respect is close to the idea of "*commensurate scales*" (*Brodsky et al*).

At the one loop level we may substitute $\alpha_{\rm eff}(\mathbf{m^2}) = \alpha_{\mathbf{s}}(\mathbf{m^2})$, develop the geometric series

$$\alpha_{\rm eff}(m^2) \simeq \alpha_s \sum_{k=0}^{\infty}\left(\frac{\beta_0\alpha_s}{4\pi}\,\ln\frac{Q^2}{m^2}\right)^k, \quad \alpha_s \equiv \alpha_s(Q^2)$$

and look for higher order perturbative corrections to our observable :

$$V(Q^2,x) - \alpha_s(Q^2)\mathcal{F}_V(0,x) \simeq \alpha_s \sum_{k=1}^{\infty}\left(\frac{\beta_0\alpha_s}{4\pi}\right)^k R_k \quad \text{with} \quad \boxed{R_k = \int_0^\infty \frac{d\epsilon}{\epsilon}\left(\ln\frac{1}{\epsilon}\right)^k \dot{\mathcal{F}}_V(\epsilon)}$$

In the IR region :

$$\dot{\mathcal{F}}_V(\epsilon) \simeq \epsilon^p\, f_V(\ln\epsilon) \implies R_k^{\rm IR} = \int_0^1 \frac{d\epsilon}{\epsilon}\left(\ln\frac{1}{\epsilon}\right)^k \epsilon^p f(\ln\epsilon) \sim \boxed{p^{-k}\,k!}$$

non-Borel--summable series !

Attempts to ascribe meaning to such a nasty series give rise to unphysical complex contributions at the level of Q^{-2p} terms : ***INFRARED RENORMALON*** problem.

This is generally interpreted as an intrinsic uncertainty in the summation of the perturbative series.

In fact, ***infrared renormalons*** are a purely ***perturbative*** phenomenon and have no direct relation to the presence of the "***Landau singularity***" in the running coupling **!**

The problem is of a physical nature and cannot be resolved by formal mathematical manipulations alone. It requires genuinely new physical input to obtain a sensible answer.

Let's introduce

$$\alpha_s(k^2) = \alpha^{\rm PT}(k^2) + \alpha^{\rm NP}(k^2)$$

It should be made clear that such a splitting is symbolic :

it represents the coupling not in terms of two ***functions*** but rather of two ***procedures***.

Having met $\alpha^{\rm PT}$ under the integral we are advised to calculate it ***perturbatively***, that is in terms of (*not too long*) a series at the point $k^2 \sim Q^2$ that our integral is "sitting" around.

At the same time we are supposed not to worry about the PT-coupling being sick in the IR region.

On the contrary, integrals with $\alpha^{\rm NP}$ are determined by that very same **IR** region and ***converge*** :

$$\int_0^\infty \frac{dk^2}{k^2}\, \alpha_s^{\rm NP}(k^2)\cdot k^{2p} \;=\; (\text{few 100s MeV})^{2p}$$

(*ITEF picture !*)

Convergence of these integrals translates into vanishing of (*first few*) integer moments of $\alpha_{\rm eff}$:

$$\lim_{k\to\infty} k^{2p}\alpha^{\rm NP}(k^2) = 0 \iff \int_0^\infty \frac{dm^2}{m^2} m^{2p}\alpha_{\rm eff}^{\rm NP}(m^2) = 0$$

(*small size instantons* $\mathbf{p} \le \beta_0 \sim 9$

thus explaining that mystery of "***non-analyticity***" in $\boldsymbol{m^2}$ *necessary to trigger on "large distances"*

The non-analyticity of ϵ, necessary to generate NP power correction, is typically of two kinds.

- In the first case an ***integer*** power $\epsilon^{\mathbf{p}}$ is accompanied by **logarithm**(**s**) of ϵ.

  This is the case of ***DIS*** structure functions, the Drell-Yan "***K-factor***", the width of hadronic ***tau-lepton*** decay, the total ***e+e- annihilation*** cross section :

$$\lim_{m \to 0} \dot{\mathcal{F}}_{\mathrm{DIS}} = a(x) \cdot \frac{m^2}{Q^2} \ln \frac{Q^2}{m^2} \quad + \ldots$$

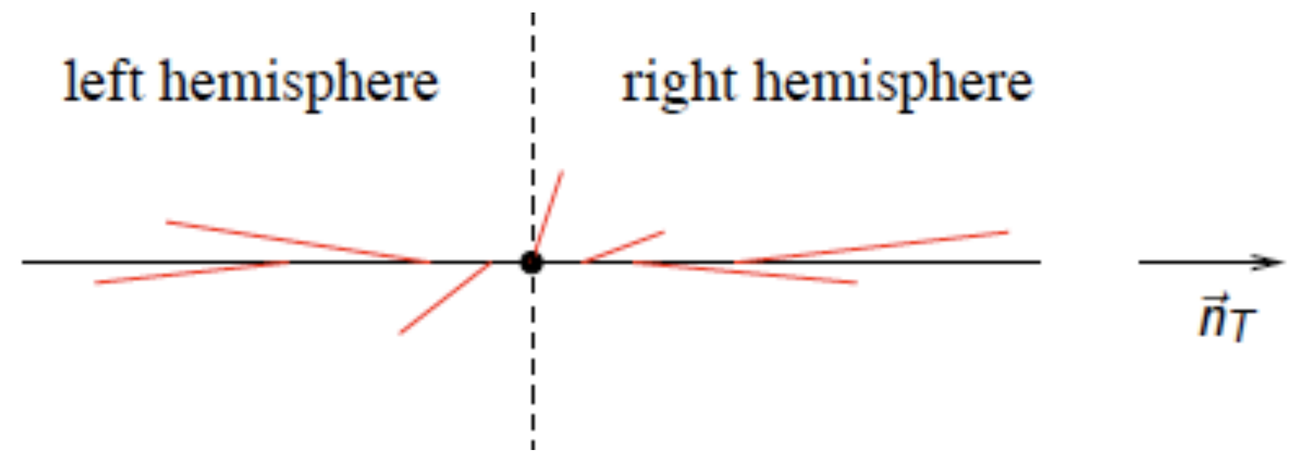
- Secondly, one may have a ***half-integer*** **p**.

  This is the case for many so-called ***jet-shape observables*** that characterise, in a CIS manner, the structure of final states produced in hard processes.

Thrust (***T***), invariant jet masses, ***C***-parameter, jet broadening (***B***), energy-energy correlation (***EEC***) etc. belong to the **p=1/2** class : they embody ***1/Q*** power effects due to confinement physics.

$$\lim_{m \to 0} \dot{\mathcal{F}}_V = a_V \cdot \sqrt{\frac{m^2}{Q^2}} \quad + \ldots$$

thrust $T = \max_{\vec{n}} \frac{\sum_i |\vec{p}_i . \vec{n}|}{\sum_i |\vec{p}_i|}$



C-param. $C = \frac{3}{2} \frac{\sum_{i,j} |\vec{p}_i||\vec{p}_j| \sin^2 \theta_{ij}}{(\sum_i |\vec{p}_i|)^2}$

jet-mass $\rho = \frac{\left(\sum_{i \in \text{hemisphere}} p_i\right)^2}{(\sum_i E_i)^2}$

broadening $B_T = \frac{\sum_i p_{ti}}{\sum_i |\vec{p}_i|}$

They are formally calculable in pQCD (being collinear and infrared safe) but possess ***large non-perturbative 1/Q–suppressed corrections.***

T is obviously a CIS observable : unaffected by either collinear parton splitting or soft gluon radiation.

Two back-to-back particles produced in the c.m.s. of ***e+e-*** annihilation correspond to ***T=1***.

Thrust deviates from unity for two reasons.

One is PT gluon bremsstrahlung : $\langle 1 - T \rangle^{(\mathrm{PT})} = \mathcal{O}(\alpha_s)$

Another reason is pure hadronization physics :

PT radiation switched off, two outgoing quarks are believed to produce two narrow jets of ***hadrons***

Hadrons are uniformly distributed in rapidity and have limited transverse momenta with respect to the jet axis (Field-Feynman hot-dog, or a Lund string).

Take a simplified "tube model" with an exponential inclusive distribution of hadrons :

$$\frac{dN}{d\eta dk_\perp} = \mu^{-1}\,\vartheta(\eta_m - |\eta|)\,e^{-k_\perp/\mu}\,, \qquad \mu = \langle k_\perp \rangle$$

Total energy :

$$\sum_i |\vec{p}_i| = 2\int d\eta\, dk_\perp \frac{dN}{d\eta dk_\perp}\, k_\perp \cosh\eta = 2\mu \sinh\eta_m = Q$$

z-momentum projections :

$$\sum_i |p_{zi}| = 2\int d\eta\, dk_\perp \frac{dN}{d\eta dk_\perp}\, k_\perp \sinh\eta = 2\mu(\cosh\eta_m - 1)$$

Constructing the ratio,

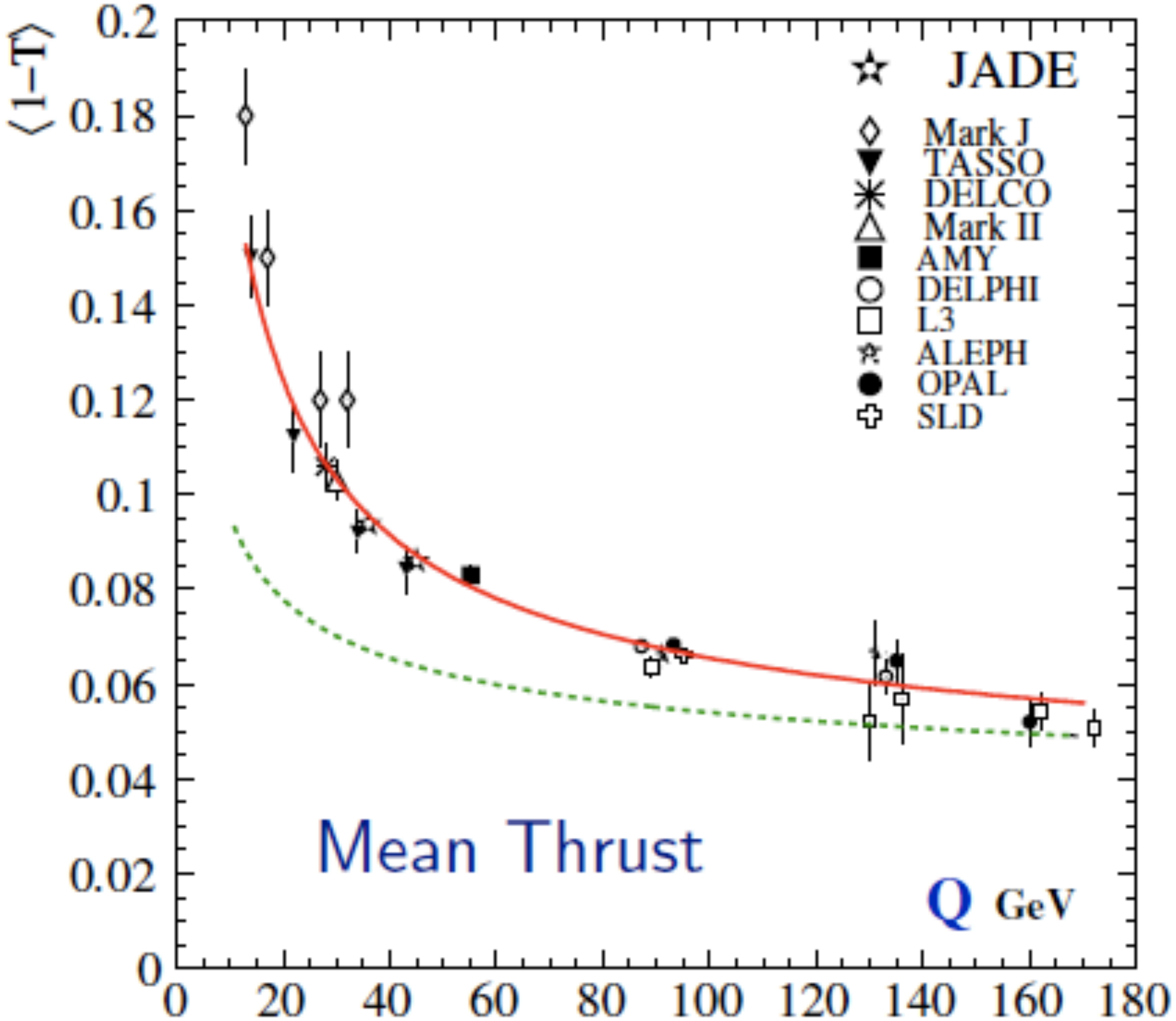
$$T = \frac{\cosh\eta_m - 1}{\sinh\eta_m} = 1 \boxed{- \frac{2\mu}{Q}} + \mathcal{O}\left(\frac{\mu^2}{Q^2}\right)$$

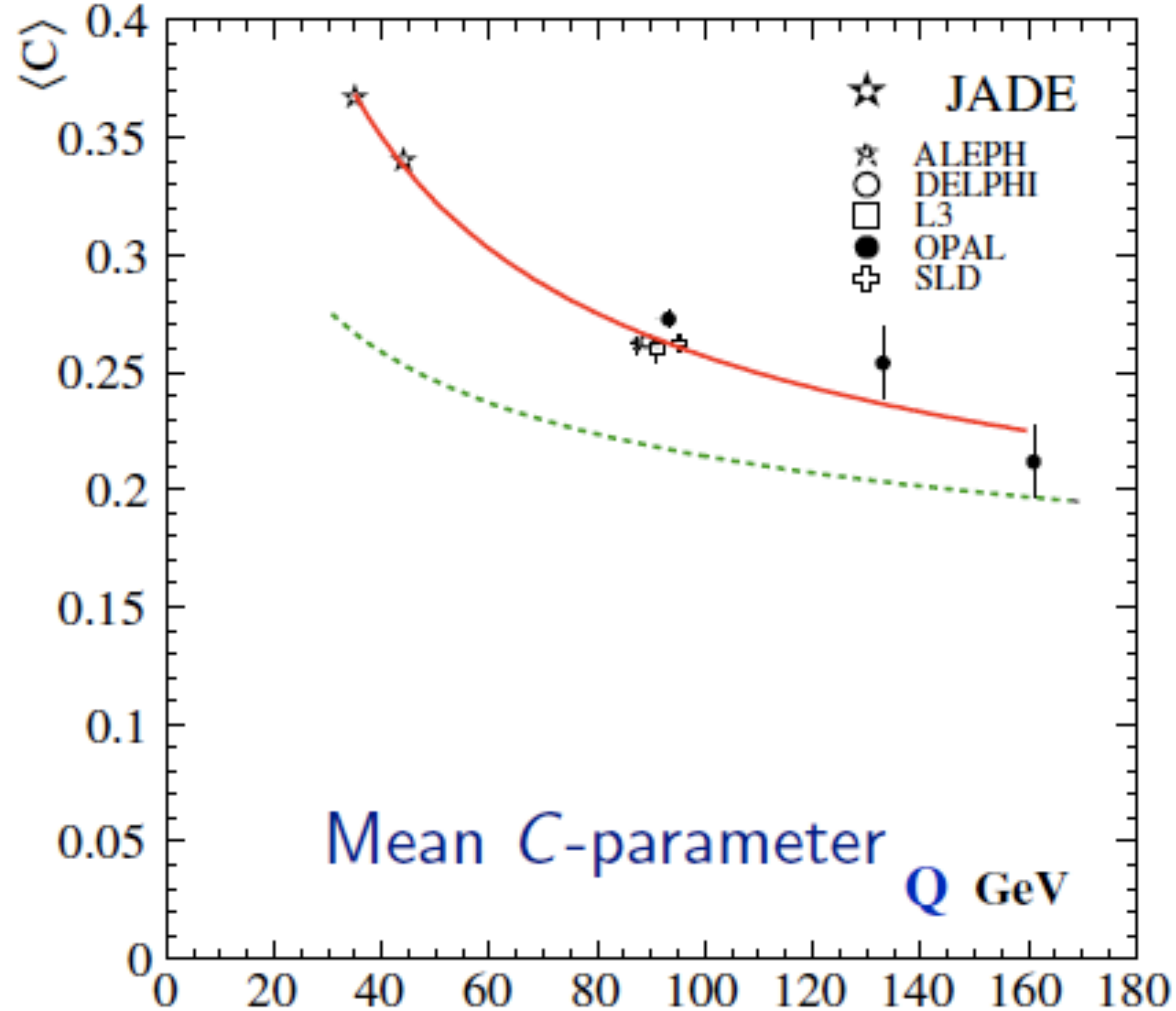
$$\boxed{\langle 1 - T\rangle^{(\mathrm{NP})} \simeq \frac{2\,\langle k_\perp\rangle}{Q}}$$

Negative ***1/Q*** ***hadronization correction !***

It is from the study of hadronization models that the ***1/Q*** effects first came into focus (*Webber, 1995*)

NP effects





"theory" :

$$\langle 1-T\rangle_{\text{hadron}} \approx \langle 1-T\rangle_{\text{parton}} + 1\ \text{GeV}/Q$$

$$\langle C\rangle_{\text{hadron}} \approx \langle C\rangle_{\text{parton}} + 4\ \text{GeV}/Q$$

$$\left\{ \frac{4}{1} \Longrightarrow \frac{3\pi}{2} \right.$$

The PT approach normally would not provide us with such a ***dimensional*** parameter: gluon transverse momenta are broadly (*logarithmically*) distributed which results in the mean $\langle k_\perp \rangle \propto \alpha_s Q$

However now we have a PT-handle on the large-distance physics : the "***gluon-mass***" trigger.

Contribution to ***thrust*** from a single gluon with momentum ***k*** reads, in terms of Sudakov variables,

$$\delta(1-T) \;=\; \min\{\alpha,\beta\} \qquad\qquad k = \alpha p + \beta \bar{p} + \mathbf{k}_\perp ,$$

$$\mathcal{F}_T \simeq \frac{C_F}{\pi}\int \frac{d\alpha}{\alpha}\frac{d\beta}{\beta} dk_\perp^2\, \delta(\alpha\beta Q^2 - k_\perp^2 - m^2)\cdot \min\{\alpha,\beta\} \;=\; \frac{2C_F}{\pi Q}\int_0^{Q^2} \frac{dk_\perp^2}{\sqrt{k_\perp^2+m^2}}$$

Evaluating the logarithmic derivative of the characteristic function,

$$\dot{\mathcal{F}} \equiv -m^2 \frac{d\mathcal{F}}{dm^2} \simeq \frac{2C_F}{\pi}\,\frac{m}{Q}$$

Thus, for non-perturbative correction to ***mean values*** of jet shape observables we get the integral

$$\langle V \rangle^{\mathrm{NP}} = \frac{a_V}{Q}\cdot \frac{C_F}{2\pi}\int_0^\infty \frac{dm^2}{m^2}\sqrt{m^2}\,\alpha_{\mathrm{eff}}^{\mathrm{NP}}(m^2)$$

where the coefficients $\mathbf{a_V}$ are simple numbers having a clear geometric origin.

The magnitude of the first two moments (jet shapes and DIS) :

$$\mathcal{A}_1 \equiv \frac{C_F}{2\pi}\int_0^\infty \frac{dk^2}{k^2}\cdot k\,\alpha_s^{\mathrm{NP}}(k^2) \;\simeq\; 0.2\text{–}0.25\,\mathrm{GeV} \qquad \mathcal{A}_2 \equiv \frac{C_F}{2\pi}\int_0^\infty \frac{dk^2}{k^2}\cdot k^2\,\alpha_s^{\mathrm{NP}}(k^2) \;\sim\; 0.2\,\mathrm{GeV}^2$$

Parametrization of the answer in terms of the ***full coupling*** *:*

$$\int_0^{\mu_I} dk\,\alpha^{\mathrm{NP}}(k^2) = \int_0^{\mu_I} dk\,\alpha_s(k^2) - \int_0^{\mu_I} dk\,\alpha^{\mathrm{PT}}(k^2) \;=\; \mu_{\mathbf{I}}\cdot\left[\alpha_{\mathbf{0}} - \left(\alpha_{\mathbf{s}}(\mathbf{Q^2}) + \beta_{\mathbf{0}}\frac{\alpha_{\mathbf{s}}^{\mathbf{2}}(\mathbf{Q^2})}{\mathbf{2\pi}}\ln\frac{\mathbf{Q}}{\mu_{\mathbf{I}}} + \dots\right)\right]$$

The characteristic non-perturbative parameter - the average of the coupling over the IR region :

$$\boxed{\alpha_0(\mu_I) \equiv \frac{1}{\mu_I}\int_0^{\mu_I} dk\, \alpha_s(k^2)}$$

Non-perturbative corrections to ***mean values*** of jet shapes $\langle V\rangle = \langle V\rangle^{\mathrm{PT}}(\alpha_s) + a_V \cdot \mathcal{P}$

with $$\mathcal{P} = \frac{4C_F\mathcal{M}}{\pi^2}\frac{\mu_I}{Q}\cdot\left\{\alpha_0(\mu_I) - \left[\alpha_s + \beta_0\frac{\alpha_s^2}{2\pi}\left(\ln\frac{Q}{\mu_I}+1+\frac{K}{\beta_0}\right)+\ldots\right]\right\}$$

the so-called "*Milan factor*" takes care of next-to-leading PT effects in the leading NP power correction

Interestingly, the same NP parameter enters the ***differential distributions*** of jet shapes :

$$\frac{d\sigma}{dV}(V) = \frac{d\sigma^{(\mathrm{PT})}}{dV}(V - a_V\mathcal{P})$$

Perturbatively calculable "*geometrical*" ***coefficients*** entering the jet shapes :

$V =$	$1-T$	C	M_T^2	M_H^2
$a_V =$	2	3π	2	1

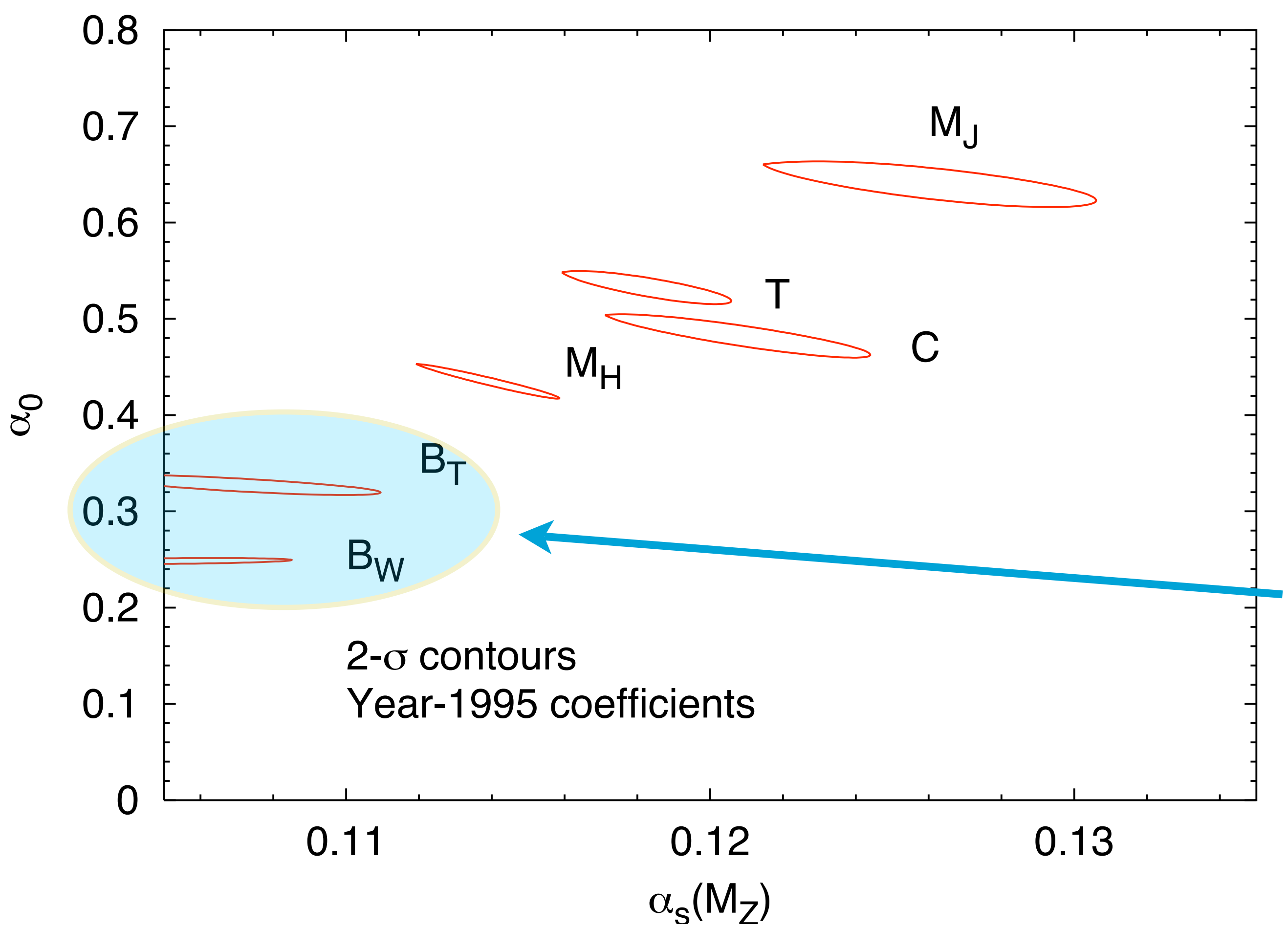
broadening drama

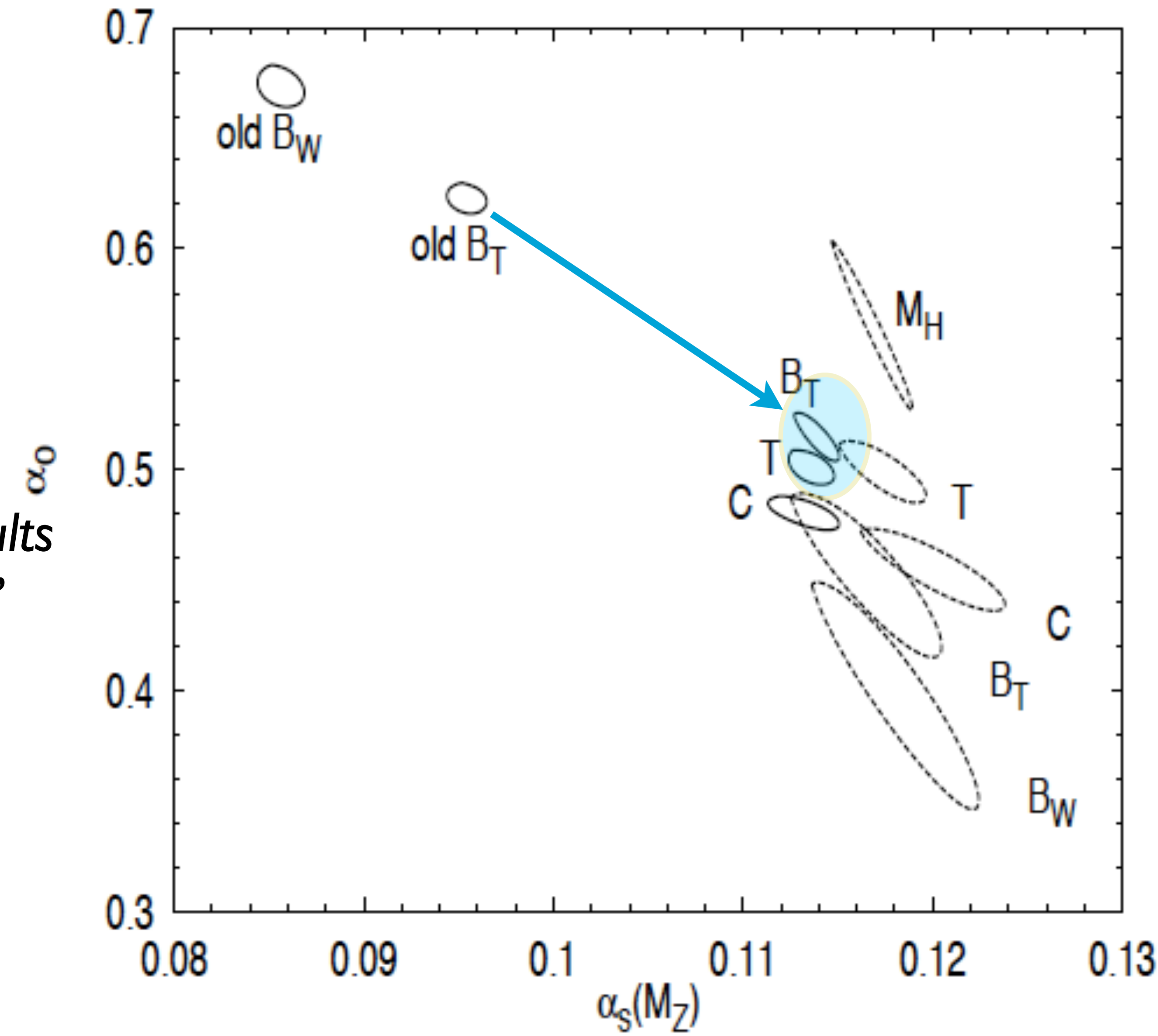
The phenomenology of power-suppressed contributions to jet shapes had a troubled childhood.
Only thrust and C-parameter remained unaffected by theoretical misconceptions...

By 1997 a ball-park value of α_0 ~0.5 was repeatedly emerging from the analyses of jet shapes in e+e- and DIS current jets in the Breit frame.

A typical resume ran like : "*The concept of a universal Power Correction parameter α_0 in DIS **ep** scattering and **e+e-** annihilation is supported*".

Montpellier 1998 QCD conference : "*inconsistent results for the total and wide-jet **broadening distributions***"





Correspondingly, the ***mean values*** of Broadenings (***T***otal and ***W***ide-jet) were found off-mark too...

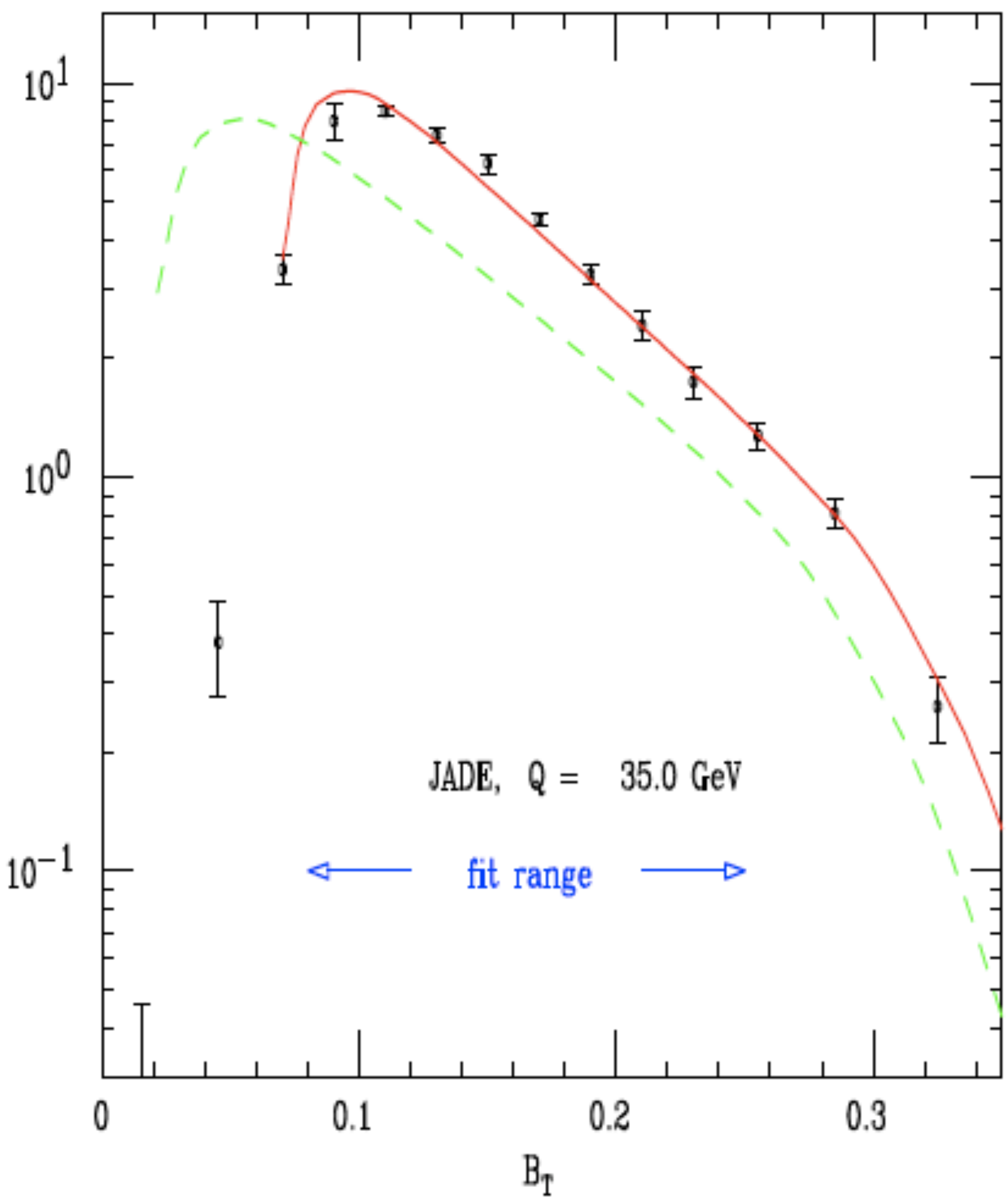
Theoretical revisiting of the Broadening measure taught an ***important lesson*** and has finally calmed the game ...

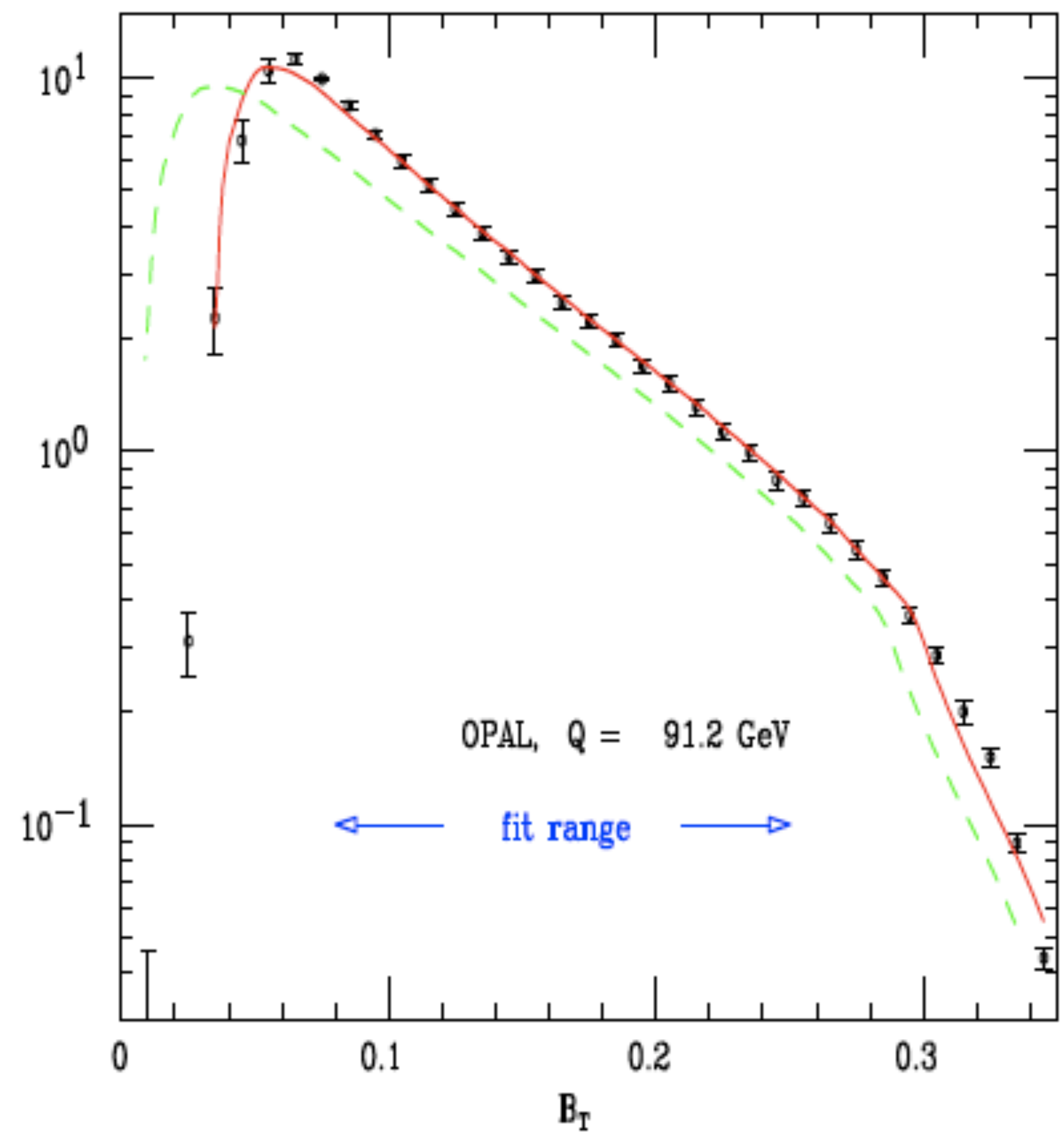
The broadening was put under scrutiny by the resurrected **JADE** collaboration.

Not only have they observed the discrepancy, but also have clarified what was going on !

They showed that hadronization effects in broadening not only ***shift*** the distribution to larger ***B*** values (as it is the case for ***1-T*** and ***C***) but also ***squeeze*** it. A bizarre observation !..

How can it be that when you ***smear*** the distribution
(moving from ***partons*** to ***hadrons***) it actually becomes ***sharper*** !?





the broadening escape

The ***B*** distribution was found to have a rich structure exhibiting ***lnB/Q*** and **lnQ/Q** NP effects.

It was soon realized that one essential phenomenon was overlooked in the original NP-treatment of broadening, namely an ***interplay*** between ***NP*** and ***PT*** phenomena.

The effects produced by gluers in the presence of normal PT gluons are different from the effects of NP-radiation inferred from a pure first-order analysis, when the PT-radiation is "switched off".

The simplest example : the story of the ***Jet Mass*** observables.

To trigger the NP-contribution we are advised to add to the parton system a ***soft gluer***.

When do so at the Born level, that is add a ***gluer*** to the quark-antiquark system as the third and only secondary parton, we find a 1/Q confinement contribution to the squared mass of the ***quark-gluer*** system - the "***heavy jet***". Meanwhile, the opposite "***lighter***" jet containing a lonely quark gets none :

$$\boxed{a_T = a_{M_T^2} = a_{M_H^2}\,, \quad a_{M_L^2} = 0}$$

$$\boxed{a_T = a_{M_T^2} = 2a_{M_H^2} = 2a_{M_L^2}}$$

There are always normal PT gluons in the game which are responsible for the bulk of the jet mass : it's not gluer's business to decide which jet is going to be *heavier*. Confinement effects are shared ***equally***.

Now we are ready to address the ***squeezed broadening*** issue.

The feature that ***1-T*** and ***C*** have in common is that the dominant NP-contribution is determined by radiation of gluers at ***large angles***. This radiation is insensitive to the tiny mismatch $\Theta_q = \mathcal{O}(\alpha_s)$ between the quark and thrust axis directions which is due to PT gluon radiation.

Therefore the ***quark momentum direction*** can be identified with the ***thrust axis***.

the broadening escape

The broadening, ***on the contrary***, accumulates contributions which do not depend on rapidity, so that the ***mismatch*** between the ***quark*** and the ***thrust*** axis matters both in the B-means and distributions.

Having naively assumed that the quark direction coincides with that of the thrust axis, ***B*** accumulated NP-contributions from ***gluers*** with rapidities up to $\eta_i \leq \eta_{\max} \simeq \ln(Q/k_{ti})$.

In this case the shift in the B-spectrum would be logarithmically enhanced, $\boxed{\Delta_B = a_B \mathcal{P} \cdot \ln \frac{Q}{Q_B}}$

High-energy gluers are collinear to the ***quark*** rather than to the thrust axis and do not contribute to B.

$$\delta B_1^{(\mathrm{NP})} \simeq a_1 \mathcal{P} \left\langle \ln \frac{1}{\Theta_q} \right\rangle$$

As a result, the NP correction to ***B*** comes out proportional to the ***quark rapidity*** !

For ***mean values*** of ***B*** observables this yields

It is the quark Sudakov form factor that describes the distribution of relative ***quark - jet axis*** angles.

$$\left\langle \ln \frac{1}{\Theta_q} \right\rangle \simeq \frac{\pi}{2\sqrt{C_F \alpha_s(Q)}}$$

How about the ***distributions*** in ***B*** ***?***

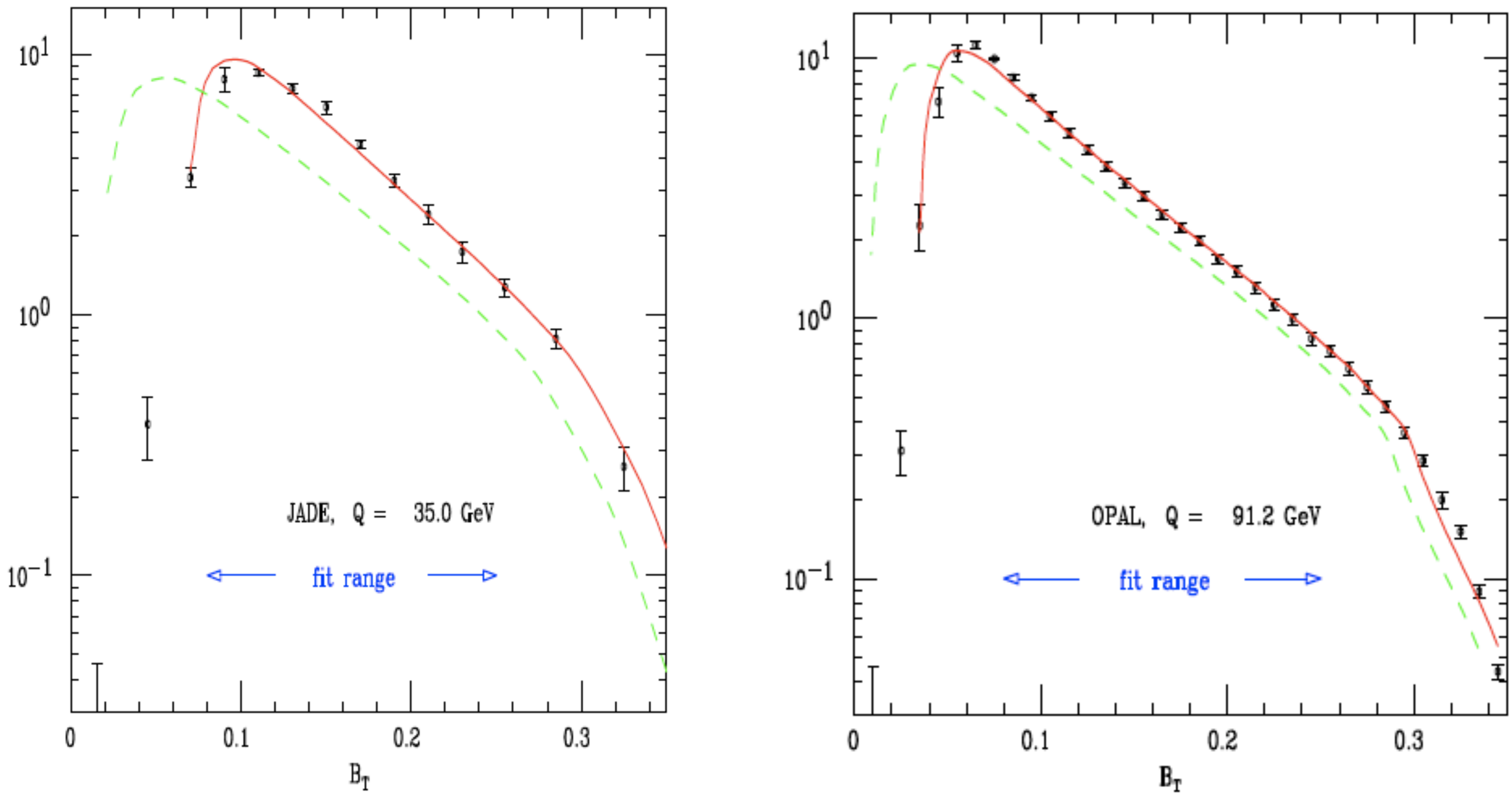
The shift in the ***single jet*** (*wide jet*) broadening is evaluated by averaging over the perturbative distribution in the ***quark angle*** while keeping the value of ***B*** fixed.

Since Θ_q is kinematically proportional to ***B*** $\Longrightarrow$ $\boxed{\Delta_1(B) \simeq a_1 \mathcal{P} \cdot \ln \frac{B_0}{B}}$

(the ***B_T*** *distribution has a somewhat more intricate structure ...)*

The smaller is ***B***, the larger the non-perturbative shift : ***squeezing***

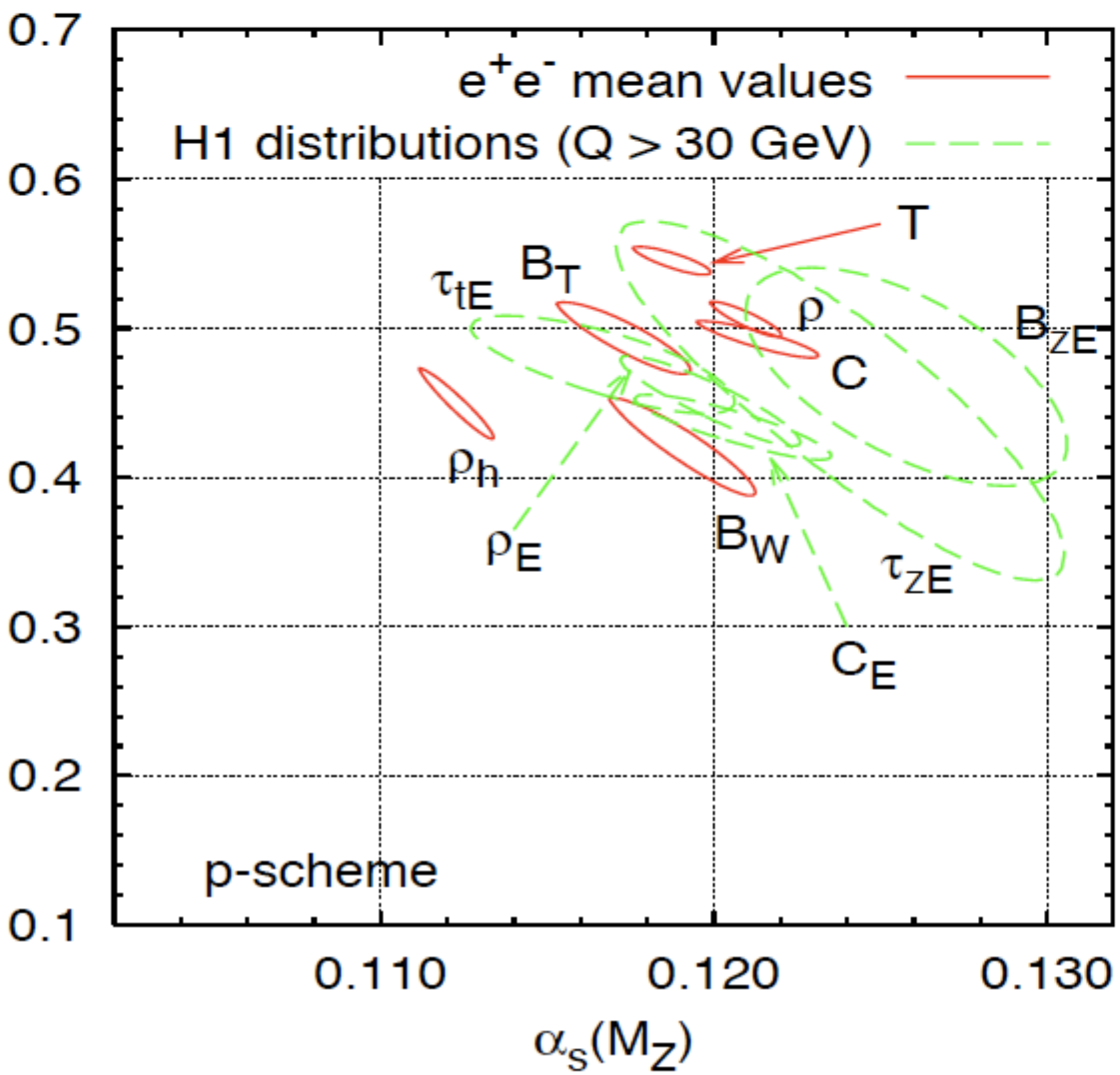
The hadronization effects in broadening not only ***shift*** the distribution to larger ***B*** values (as it is the case for ***1-T*** and ***C***) but also ***squeeze*** it.

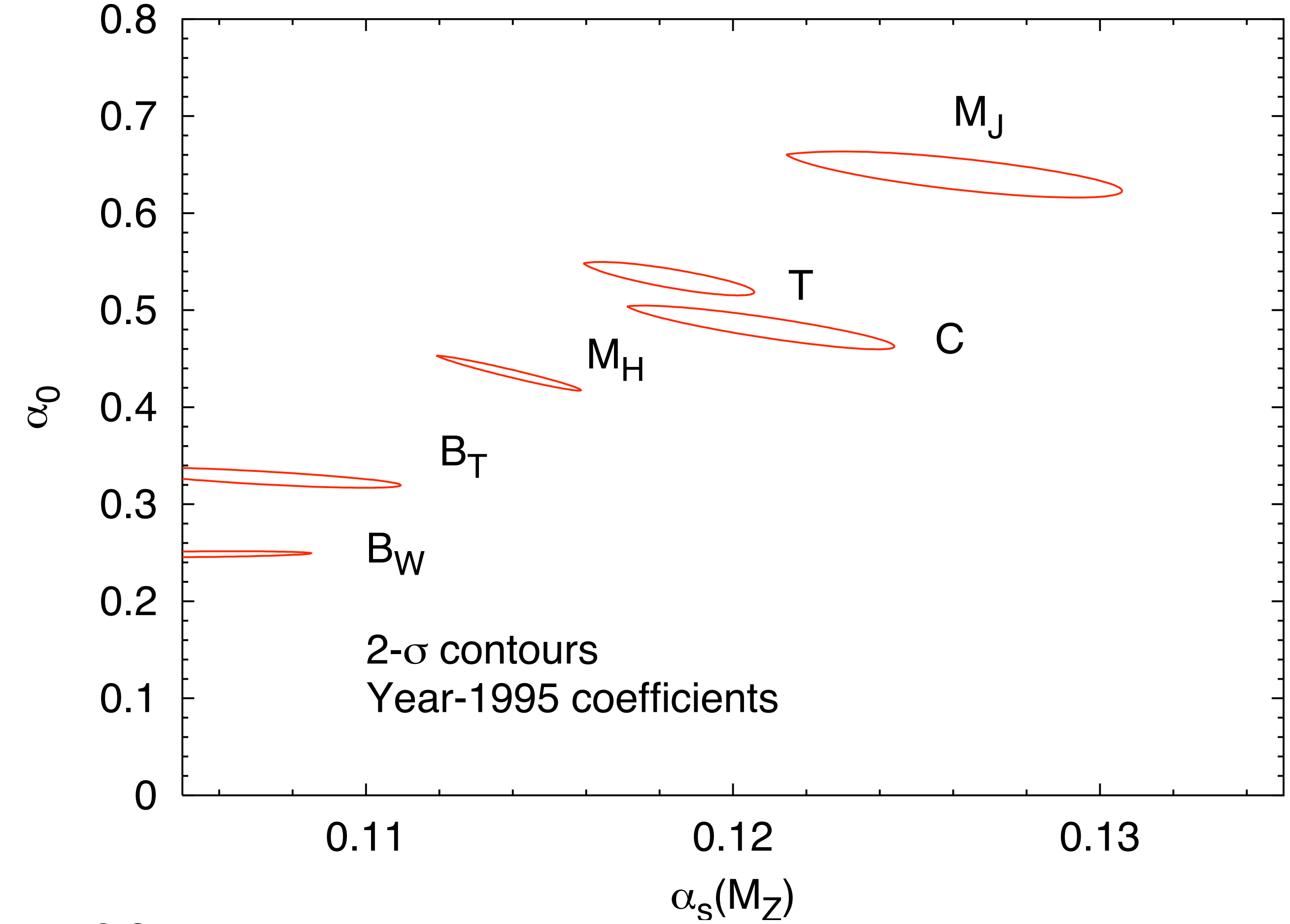


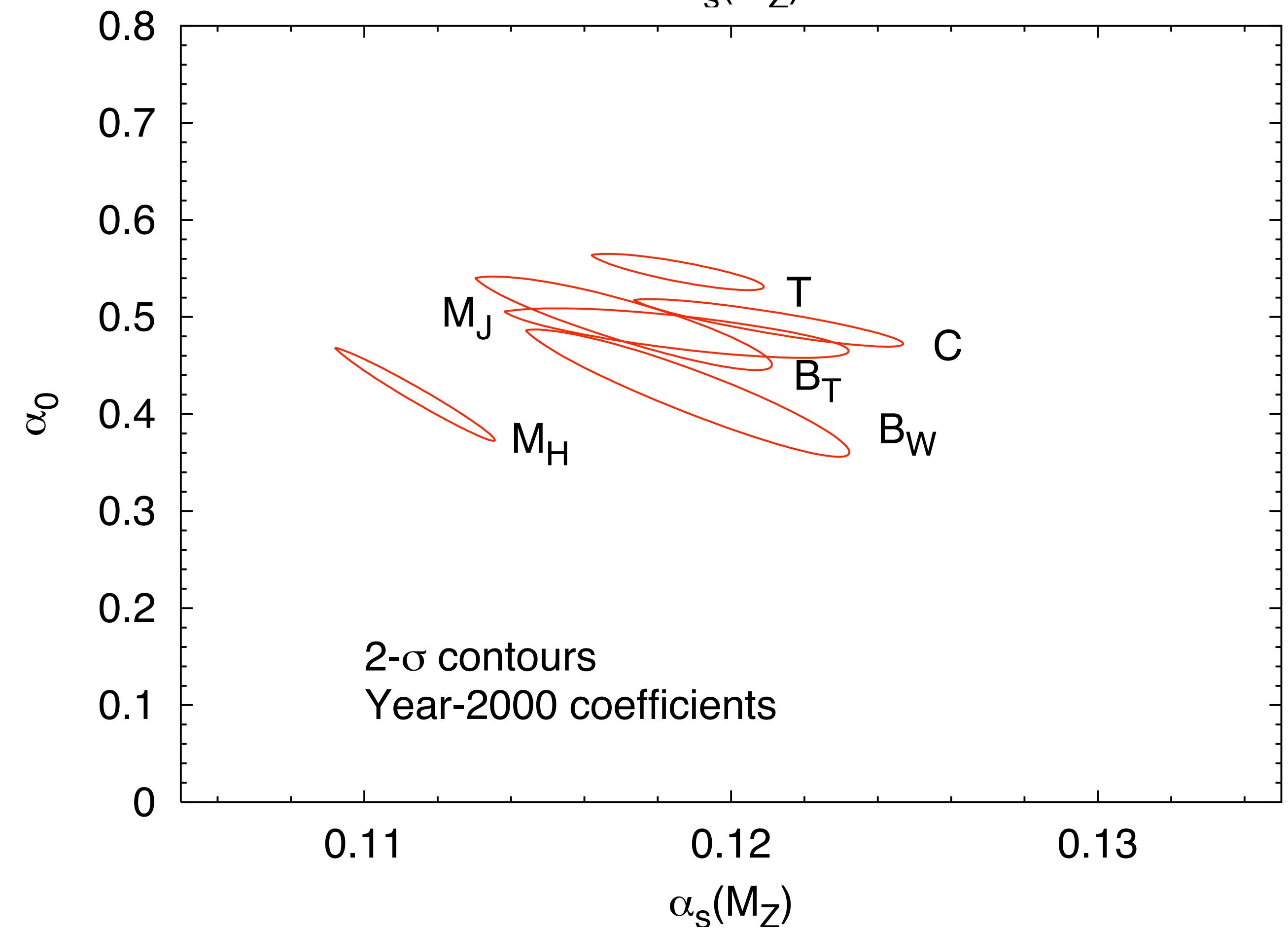
NP effects in jet shapes

After the mis-concepts and errors have been fixed, the ensemble of jet shape measurements clustered :

Moreover, both ***LEP*** (+ JADE) and ***HERA*** have contributed to the scrutiny and to the convergence







Theory + Phenomenology of 1/Q effects in event shape observables, both in e+e– annihilation and DIS systematically pointed at the average value of the infrared coupling

$$\alpha_0 \equiv \frac{1}{2\ \mathrm{GeV}} \int_0^{2\ \mathrm{GeV}} dk\, \alpha_s(k^2) \sim 0.5$$

The main features of this result are as follows : the average IR coupling is

- Universal holds to within $\pm 15\%$

  If not for the ***universality***,
  the whole game would made no sense : it would have meant just trading ***one unknown*** - non-perturbative "smearing" effects in a given observable (like in MC event generators) - for ***another unknown*** function - the shape of the coupling in the infrared...

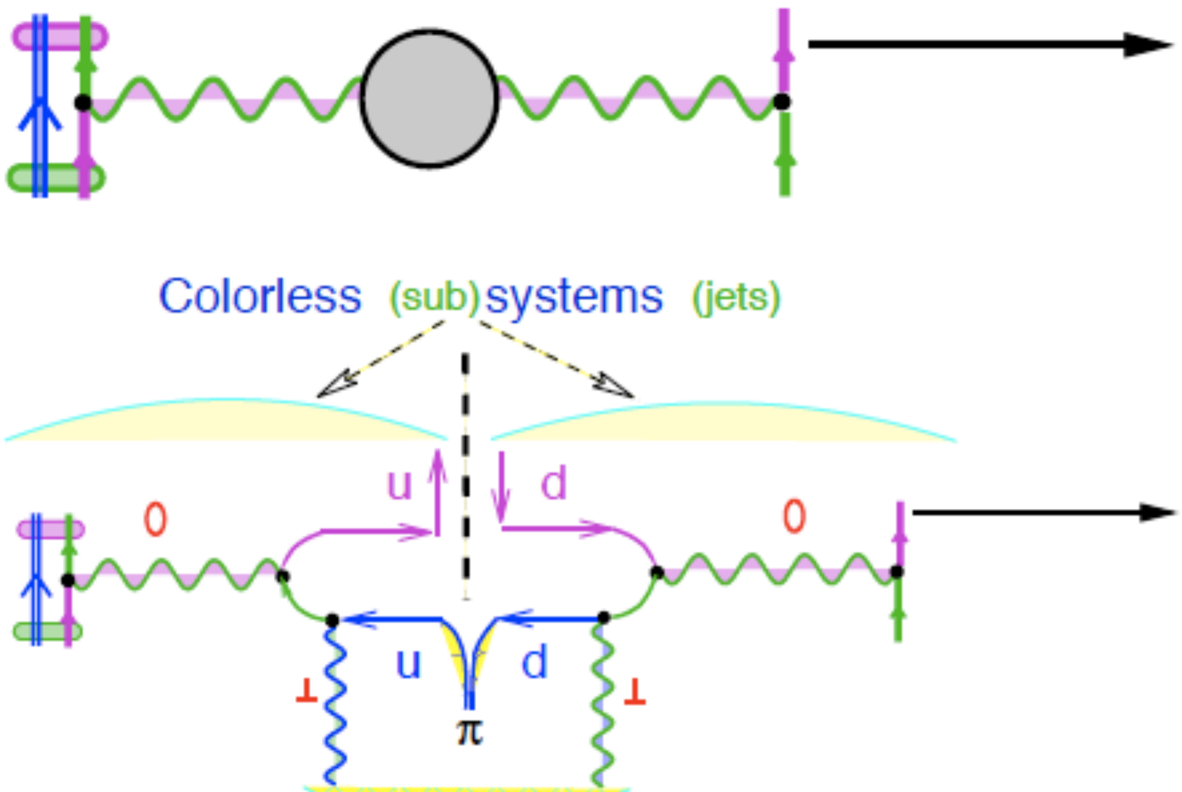
- Reasonably small (which opens intriguing possibilities . . .)

- Comfortably above the Gribov's critical value ($\pi \cdot 0.137 \simeq 0.4$)

EXTRAS

A word about "*soft confinement*"

What happens with the Coulomb field when the sources move apart?



Bearing in mind that virtual quarks live in the background of gluons (zero fluctuations of $A_\perp$ gluon fields) what we look for is a mechanism for binding (negative energy) *vacuum quarks* into *colorless hadrons* (positive energy physical states of the theory)

V.Gribov suggested such a mechanism — the supercritical binding of light fermions subject to a Coulomb-like interaction. It develops when the coupling constant hits a definite "critical value" (Gribov 1990)

$$\boxed{\frac{\alpha}{\pi} > \frac{\alpha_{\rm crit}}{\pi} = C_F^{-1}\left[1 - \sqrt{\frac{2}{3}}\right] \simeq 0.137}$$

$$\left(C_F = \frac{N_c^2-1}{2N_c} = \frac{4}{3}\right)$$

An amazing success of the relativistic theory of electron and photon fields — quantum electrodynamics (QED) — has produced a long-lasting negative impact: it taught the generations of physicists that came into the business in/after the 70's to "not to worry".

Indeed, today one takes a lot of things for granted :

- One rarely questions whether the alternative roads to constructing QFT — *secondary quantization*, *functional integral* and the *Feynman diagram* approach — really lead to the ***same quantum theory*** of interacting fields

- One feels ashamed to doubt an elegant powerful, *but potentially deceiving,* technology of translating the dynamics of quantum fields into that of statistical systems

- One takes the original concept of the "***Dirac sea*** " — the picture of the fermionic content of the vacuum — as an anachronistic model

- One was taught to look upon the problems that arise with field-theoretical description of point-like objects and their interactions at very small distances (*ultraviolet divergences*) as purely technical : ***renormalize it and forget it***.

QED : physical objets — ***electrons*** and ***photons*** — are in *one-to-one correspondence* with the fundamental fields that one puts into the local Lagrangian of the theory.

The rˆole of the QED Vacuum is "trivial": it makes e.m. charge (and the electron mass operator) run, but does not affect the nature of the interacting fields.

QCD : the Vacuum changes the bare fields ***beyond recognition*** ...

Gribov Confinement: setting up the Problem

- The question of interest is

  the confinement in real world (with 2 very light ***u*** and ***d*** quarks), rather than ***a*** confinement.

- No mechanism for binding massless bosons (gluons) seems to exist in Quantum Field Theory (QFT), while the Pauli exclusion principle may provide means for binding together massless fermions (light quarks).

- The problem of ultraviolet regularization may be more than a technical trick in a QFT with apparently ***infrared-unstable dynamics*** : the ***ultraviolet*** and ***infrared*** regimes of the theory may be tightly linked.

- The Feynman diagram technique has to be reconsidered in QCD if one goes beyond trivial perturbative correction effects.

  Feynman's famous $i\epsilon$ prescription was designed for (*and applies only to*) quantum field theories with ***stable perturbative vacua***.

- To understand and describe a physical process in a ***confining theory***, it is necessary to take into consideration the ***response of the vacuum***, which leads to essential modifications of the quark and gluon Green functions.

A known QFT example of such a violent response of the vacuum — screening of ***super-charged ions*** with ***Z > 137***.

The expression for Dirac energy levels of an electron in a field created by the point-like electric charge Z contains $\epsilon \propto \sqrt{1-(\alpha_{\rm e.m.}Z)^2}$.

For $Z > 137$ the energy becomes *complex*. This means instability.

- Classically, the electron "*falls onto the centre*".
- Quantum-mechanically, it also "*falls*", but into the Dirac sea.

$$A_Z \Longrightarrow A_{Z-1} + e^+ , \qquad \text{for } Z > Z_{\rm crit.}$$

(Pomeranchuk & Smorodinsky 1945)

In the QCD context, the increase of the running quark-gluon coupling at large distances replaces the large Z of the QED problem.

Gribov generalized the problem of supercritical binding in the field of an ***infinitely heavy source*** to the case of two ***massless fermions*** interacting via Coulomb-like exchange.

He found that in this case the supercritical phenomenon develops much earlier.

Namely, a pair of light fermions develops supercritical behaviour if the coupling hits a definite critical value

$$\frac{\alpha}{\pi} > \frac{\alpha_{\rm crit}}{\pi} = 1 - \sqrt{\frac{2}{3}}$$

With account of the QCD color Casimir operator, the value of the coupling above which restructuring of the perturbative vacuum leads to ***chiral symmetry breaking*** and, *likely*, to ***confinement*** , translates into

$$\boxed{\frac{\alpha_{\rm crit}}{\pi} = C_F^{-1}\left[1 - \sqrt{\frac{2}{3}}\right] \simeq 0.137}$$