QCD and Hadron Dynamics

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1. Introduction

Quantum Chromodynamics is a wonder. Take free quarks, supply them with colour degrees of freedom, demand invariance with respect to the "repainting" quark fields arbitrarily in each point in space-time — and you get the unique QCD Lagrangian describing interacting quarks and gluons. Having done that, you (are supposed to) have the whole hadron world in your hands! Such a beauty and ambition is hard to match.

At the same time, it is worth remembering that QCD is probably the strangest of theories in the history of modern physics. On one hand, the striking successes of QCD-based phenomenology leave no doubt that QCD is indeed the microscopic theory of hadrons and their interactions. On the other hand, the depth of the conceptual problems that one faces in trying to formulate QCD as a respectable Quantum Field Theory is unprecedented.

QCD nowadays has a split personality. It embodies "hard" and "soft" physics, both being hard subjects, and the softer the harder.

Until recently QCD studies were concentrated on small-distance phenomena, observables and characteristics that are as insensitive to large-distance confinement physics as possible. This is the realm of "hard processes" in which a large momentum transfer Q^2, either time-like $Q^2 \gg 1$ GeV^2, or space-like $Q^2 \ll -1$ GeV^2, is applied to hadrons in order to probe their small-distance quark-gluon structure.

High-energy annihilation $e^+e^- \to$ hadrons, deep inelastic lepton-hadron scattering (DIS), production in hadron-hadron collisions of massive lepton pairs, heavy quarks and their bound states, large transverse momentum jets and photons are classical examples of hard processes.

Perturbative QCD (PT QCD) controls the relevant cross sections and, to a lesser extent, the structure of final states produced in hard interactions. Whatever the hardness of the process, it is hadrons, not quarks and gluons, that hit the detectors. For this reason alone, the applicability of the PT QCD approach, even to hard processes, is far from being obvious. One has to rely on plausible arguments (completeness, duality) and look for observables that are less vulnerable towards our ignorance about confinement.

In particle physics a discovery of a class of animals that are *more equal* than the others is due to Sterman & Weinberg (1977). They introduced an important notion of Collinear-and-Infrared Safety.

An observable is granted the CIS status if it can be calculated in terms of quarks and gluons treated as real particles (partons), without encountering either collinear ($\theta \to 0$) or infrared ($k_0 \to 0$) divergences. The former divergence is a standard feature of (massless) QFT with dimensionless coupling, the latter is typical for massless vector bosons (photons, gluons).

This classification is more than mere zoology. Given CIS quantity, we expect its PT QCD value *predictable* in the quark-gluon framework to be directly comparable with its *measurable* value in the hadronic world. For this reason the CIS observables are the preferred pets of QCD practitioners.

To give an example, we cannot deduce from the first principles parton distributions inside hadrons (PDF, or structure functions). However, the rate of their $\ln Q^2$-dependence (scaling violation) is an example of a CIS measure and stays under PT QCD jurisdiction.

Speaking about the final state structure, we cannot predict, say, the kaon multiplicity or the pion energy spectrum. However, one can decide to be not too picky and concentrate on global characteristics of the final states rather than on the yield of specific hadrons. Being sufficiently inclusive with respect to final hadron species, one can rely on a picture of the energy-momentum flow in hard collisions supplied by PT QCD — the jet pattern.

There are well elaborated procedures for counting jets (CIS jet finding algorithms) and for quantifying the internal structure of jets (CIS jet shape variables). They allow the study of the gross features of the final states while staying away from the physics of hadronisation. Along these lines one visualizes asymptotic freedom, checks out gluon spin and colour, predicts and verifies scaling violation pattern in hard cross sections, etc. These and similar checks have constituted the basic QCD tests of the past two decades.

This epoch is over. Now the High Energy Particle physics community is trying to probe genuine confinement effects in hard processes to learn more about strong interactions. The programme is ambitious and provocative. Friendly phenomenology keeps it afloat and feeds our hopes of extracting valuable information about physics of hadronisation.

The rôle of HERA in this quest is difficult to overestimate. She has already taught us a lot about the structure of the proton, its quark and gluon content, and brought back into limelight, in quite a dramatic fashion, such basic issues as the Pomeron, diffraction, unitarity.

It is true that e^+e^- annihilation provides the cleanest environment for studying the QCD interactions and hadron jets. The advantage of HERA, however, is that in the DIS environment the total energy of the collision and the hardness of the interaction are not strictly linked as is the case for the point-like e^+e^- annihilation.

The power of HERA, the DIS goddess, lies in her ability to study various scales, from very large Q^2 down to moderate and small momentum transfers, thus probing an interface between hard and soft physics. While it can be said that e^+e^- will always remain the best ground for *testing* QCD, DIS is better suited for *understanding* it.

2. Multihadron production and QCD

In general, there are three ways to probe the small-distance hadron structure.

vac→hadrons: High energy vacuum excitation producing hadrons, like in (but not exclusively) e^+e^- annihilation.

vac+h→hadrons: Large momentum transfer excitation of an initial hadron by a sterile probe, like in deep inelastic lepton-hadron scattering (DIS).

h+h→hadrons: Production of large–$p_\perp$ hadrons in hadron-hadron collisions. (Here sterile probes can be employed in the final state as well, e.g. massive lepton pairs and/or large–$p_\perp$ photons.)

Copious production of hadrons is typical for all these processes. On the other hand, at the microscopic level, multiple quark-gluon "production" is to be expected as a result of QCD bremsstrahlung — gluon radiation accompanying abrupt creation/scattering of colour partons.

Is there a correspondence between observable hadron and calculable quark-gluon production?

An indirect evidence that gluons are there, and that they behave, can be obtained from the study of the scaling violation pattern. QCD quarks (and gluons) are not point-like particles, as the orthodox parton model once assumed. Each of them is surrounded by a proper field coat — a coherent virtual cloud consisting of gluons and "sea" $q\bar{q}$ pairs. A hard probe applied to such a dressed parton breaks coherence of the cloud. Constituents of these field fluctuations are then released as particles accompanying the hard interaction. The harder the hit, the larger an intensity of bremsstrahlung and, therefore, the fraction of the energy-momentum of the dressed parton that the bremsstrahlung quanta typically carry away. Thus we should expect, in particular, that the probability that a "bare" core quark carries a large fraction of the energy of its dressed parent will decrease with increase of Q^2. And so it does.

The logarithmic scaling violation pattern in DIS structure functions is well established and meticulously follows the QCD prediction based on the parton evolution picture.

In DIS we look for a "bare" quark inside a target dressed one. In e^+e^- hadron annihilation at large energy $s = Q^2$ the chain of events is reversed. Here we produce instead a bare quark with energy $Q/2$, which then "dresses up". In the process of restoring its proper field-coat our parton produces (a controllable amount of) bremsstrahlung radiation which leads to formation of a hadron jet. Having done so, in the end of the day it becomes a constituent of one of the hadrons that hit the detector. Typically, this is the leading hadron. However, the fraction x_E of the initial energy $Q/2$ that is left to the leader depends on the amount of accompanying radiation and, therefore, on Q^2 (the larger, the smaller). In fact, the same rule (and the same formula) applies to the scaling violation pattern in e^+e^- fragmentation functions (time-like parton evolution) as to that in the DIS parton distributions (space-like evolution).

What makes the annihilation channel particularly interesting, is that the present day experiments are so sophisticated that they provide us with a near-to-perfect separation between quark- and gluon-initiated jets (the latter being extracted from heavy-quark-tagged three-jet events).

In Fig. 1 a comparison is shown of the scaling violation rates in the hadron spectra from gluon and quark jets, as a function of the hardness scale κ that char-

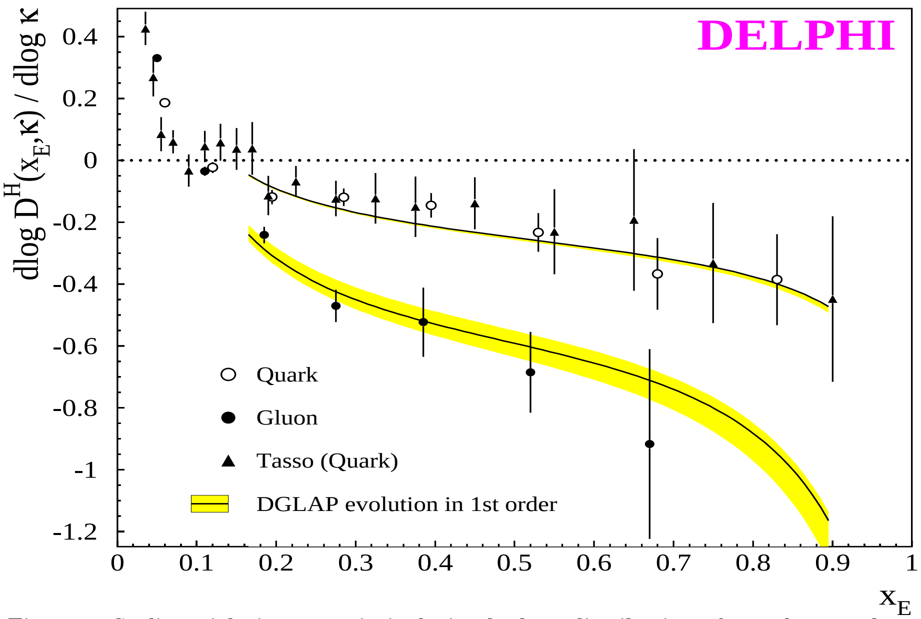

Figure 1. Scaling violation rates in inclusive hadron distributions from gluon and quark jets (Hamacher *et. al* 1999).

acterizes a given jet (Hamacher *et al.* 1999). For large values of $x_E \sim 1$ the ratio of the logarithmic derivatives is predicted to be close to that of the gluon and quark "colour charges", $C_A/C_F = 9/4$. Experimentally, the ratio is measured to be

$$\frac{C_A}{C_F} = 2.23 \pm 0.09_{\mathrm{stat.}} \pm 0.06_{\mathrm{syst.}}. \tag{2.1}$$

(*a*) *Mean parton and hadron multiplicities*

Since accompanying QCD radiation seems to be there, we can make a step forward by asking for a *direct* evidence: what is the fate of those gluons and sea quark pairs produced via multiple initial gluon bremsstrahlung followed by parton multiplication cascades? Let us look at the Q-dependence of the mean hadron multiplicity, the quantity dominated by relatively soft particles with $x_E \ll 1$. This is the kinematical region populated by accompanying QCD radiation.

Fig. 2 demonstrates that the hadron multiplicity increases with the hardness of the jet proportional to the multiplicity of secondary gluons and sea quarks. The ratio of the slopes, once again, provides an independent measure of the ratio of the colour charges, which is consistent with (2.1) (DELPHI 1999):

$$\frac{C_A}{C_F} = 2.246 \pm 0.062_{\mathrm{stat.}} \pm 0.008_{\mathrm{syst.}} \pm 0.095_{\mathrm{theo.}}. \tag{2.2}$$

(*b*) *Inclusive hadron distribution in jets*

Since the total numbers match, it is time to ask a more delicate question about energy-momentum distribution of final hadrons versus that of the underlying parton

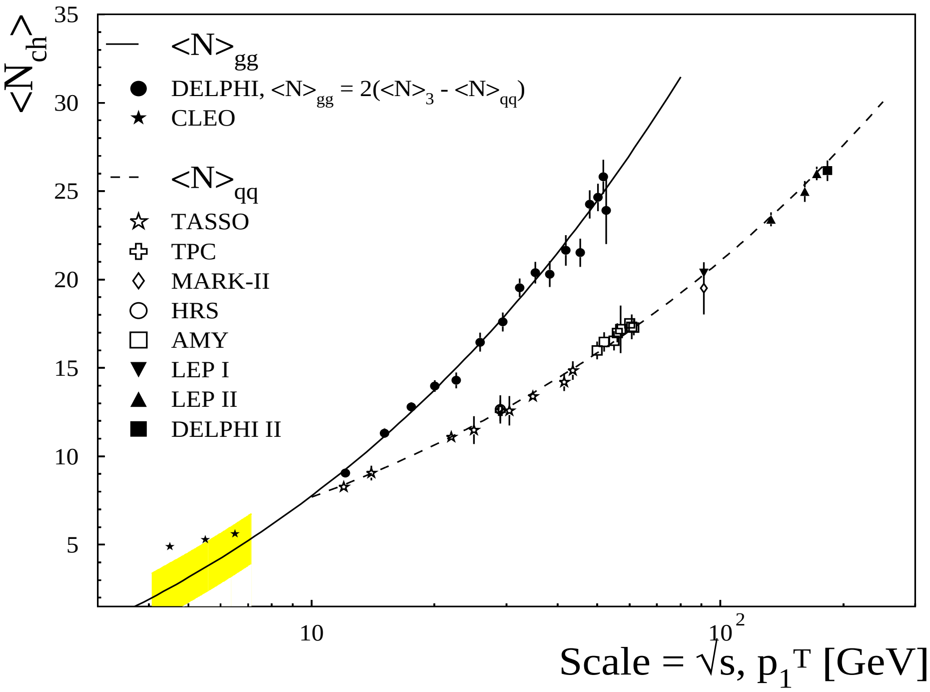

Figure 2. Charged hadron multiplicities in gluon and quark jets (DELPHI 1999).

ensemble. One should not be too picky in addressing such a question. It is clear that hadron-hadron correlations, for example, will show resonant structures about which the quark-gluon speaking PT QCD can say little, if anything, at the present state of the art. Inclusive single-particle distributions, however, have a better chance to be closely related. Triggering a single hadron in the detector, and a single parton on paper, one may compare the structure of the two distributions to learn about dynamics of hadronisation.

Inclusive energy spectrum of soft bremsstrahlung partons in QCD jets has been derived in 1984 in the so-called MLLA — the Modified Leading Logarithmic Approximation (Dokshitzer & Troyan 1984). This approximation takes into account all essential ingredients of parton multiplication in the next-to-leading order. They are: parton splitting functions responsible for the energy balance in parton splitting, the running coupling $\alpha_s(k_\perp^2)$ depending on the relative transverse momentum of the two offspring and exact angular ordering. The latter is a consequence of soft gluon coherence and plays an essential rôle in parton dynamics. In particular, gluon coherence suppresses multiple production of very small momentum gluons. It is particles with intermediate energies that multiply most efficiently. As a result, the energy spectrum of relatively soft secondary partons in jets acquires a characteristic hump-backed shape. The position of the maximum in the logarithmic variable $\xi = -\ln x$, the width of the hump and its height increase with Q^2 in a predictable way.

The shape of the inclusive spectrum of all charged hadrons (dominated by $\pi^\pm$) exhibits the same features. This comparison, pioneered by Glen Cowan (ALEPH) and the OPAL collaboration, has later become a standard test of analytic QCD predictions. First scrutinized at LEP, the similarity of parton and hadron energy distributions has been verified at SLC and KEK e^+e^- machines, as well as at HERA and Tevatron where hadron jets originate not from bare quarks dug up from the vacuum by a highly virtual photon/Z^0 but from hard partons kicked out from initial hadron(s).

In Fig. 3 (DELPHI) the comparison is made of the all-charged hadron spectra at various annihilation energies Q with the so-called "distorted Gaussian" fit (Fong & Webber 1989) which employs the first four moments (the mean, width, skewness and kurtosis) of the MLLA distribution around its maximum.

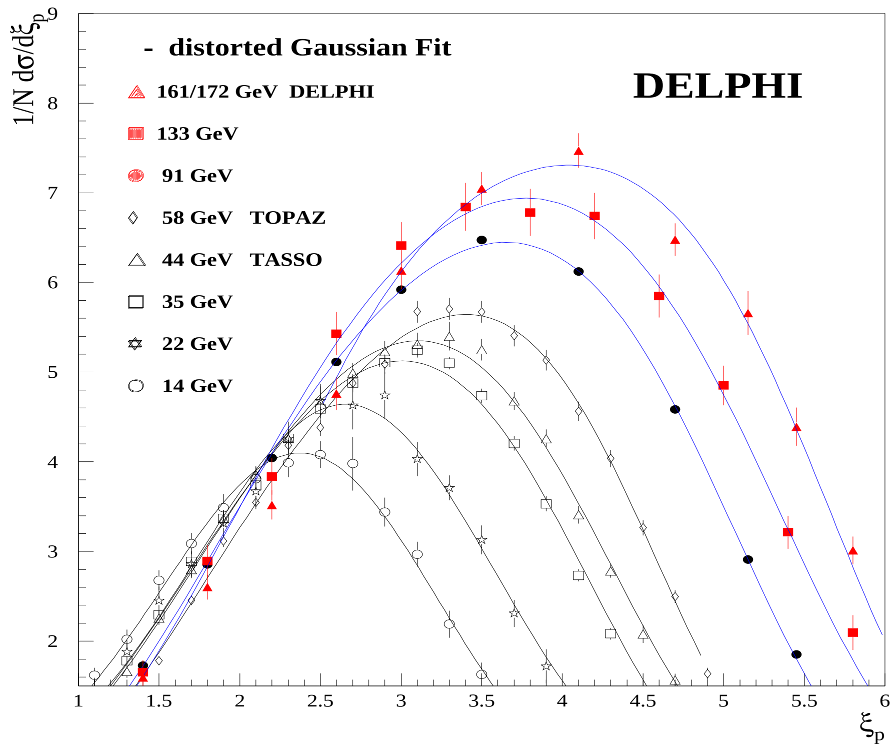

Figure 3. Inclusive energy distribution of charged hadrons in jets produced in e^+e^- annihilation

Shall we say, a (routine, interesting, wonderful) check of yet another QCD prediction? Better not. Such a close similarity offers a deep puzzle, even a worry, rather than a successful test. Indeed, after a little exercise in translating the values of the logarithmic variable $\xi = \ln(E_{\text{jet}}/p)$ in Fig. 3 into GeVs you will see that the actual hadron momenta at the maxima are, for example, $p{=}\frac{1}{2}Q \cdot e^{-\xi_{\max}} \simeq 0.42$, 0.85 and 1.0 GeV for $Q{=}14$, 35 GeV and at LEP-1, $Q{=}91$ GeV. Is it not surprising that the PT QCD spectrum is mirrored by that of the pions (which constitute 90% of all charged hadrons produced in jets) with momenta well below 1 GeV?!

For this very reason the observation of the parton-hadron similarity was initially met with a serious and well grounded scepticism: it looked more natural (and was more comfortable) to blame the finite hadron mass effects for falloff of the spectrum at large ξ (small momenta) rather than seriously believe in applicability of the PT QCD consideration down to such disturbingly small momentum scales.

This worry has been recently answered. Andrey Korytov (CDF) was the first to hear a theoretical hint and carry out a study of the energy distribution of hadrons produced inside a restricted angular cone Θ around the jet axis. Theoretically, it is not the energy of the jet but the maximal parton transverse momentum inside it, $k_{\perp\max} \simeq E_{\rm jet} \sin\frac{\Theta}{2}$, that determines the hardness scale and thus the yield and the distribution of the accompanying radiation.

This means that by choosing a small opening angle one can study relatively small hardness scales but in a cleaner environment: due to the Lorentz boost effect, eventually all particles that form a short small-Q^2 QCD "hump" are now relativistic and concentrated at the tip of the jet.

For example, selecting hadrons inside a cone $\Theta \simeq 0.14$ around an energetic quark jet with $E_{\rm jet} \simeq 100$ GeV (LEP-2) one should see that very "dubious" $Q = 14$ GeV curve in Fig. 3 but now with the maximum boosted from 450 MeV into a comfortable 6 GeV range.

In the CDF Fig. 4 (A. Korytov 1996, personal communication; Goulianos 1997, see also Safonov 1999) a close similarity between the hadron yield and the full MLLA parton spectra can no longer be considered accidental and be attributed to non-relativistic kinematical effects.

(*c*) *Brave gluon counting*

Modulo $\Lambda_{\rm QCD}$, there is only one unknown in this comparison, namely, the overall normalisation of the spectrum of hadrons relative to that of partons (bremsstrahlung gluons).

Strictly speaking, there should/could have been another free parameter, the one which quantifies one's bravery in applying the PT QCD dynamics. It is the minimal transverse momentum cutoff in parton cascades, $k_\perp > Q_0$. The strength of successive $1 \to 2$ parton splittings is proportional to $\alpha_s(k_\perp^2)$ and grows with $k_\perp$ decreasing. The necessity to terminate the process at some low transverse momentum scale where the PT coupling becomes large (and eventually hits the formal "Landau pole" at $k_\perp = \Lambda_{\rm QCD}$) seems imminent. Surprisingly enough, it is not.

Believe it or not, the inclusive parton energy distribution turns out to be a CIS QCD prediction. Its crazy $Q_0 = \Lambda_{\rm QCD}$ limit (the so-called "limiting spectrum") is shown by solid curves in Fig. 4.

Choosing the minimal value for the collinear parton cutoff Q_0 can be looked upon as shifting, as far as possible, responsibility for particle multiplication in jets to the PT dynamics. This brave choice can be said to be dictated by experiment, in a certain sense. Indeed, with increase of Q_0 the parton parton distributions *stiffen* (parton energies are limited from below by the kinematical inequality $xE_{\rm jet} \equiv k \geq k_\perp > Q_0$). The maxima would move to larger x (smaller ξ), departing from the data.

A clean test of "brave gluon counting" is provided by Fig. 5 where the position of the hump, which is insensitive to the overall normalisation, is compared with the parameter-free MLLA QCD prediction (Safonov 1999).

A formal explanation of the tolerance of the *shape* of inclusive parton spectra

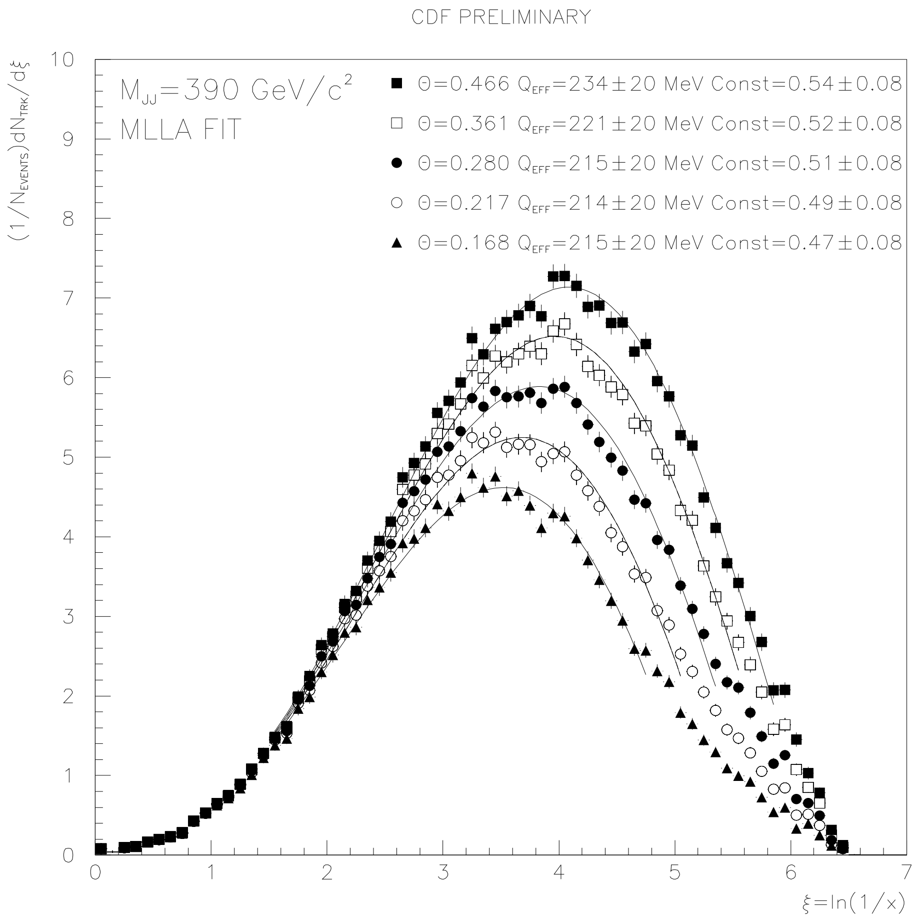

Figure 4. Inclusive energy distribution of charged hadrons in large–$p_\perp$ jets (Goulianos 1997).

to the dangerous small-$k_\perp$ domain can be found in the proceedings of the last year Blois conference (Dokshitzer 1999).

To put a long story short, decreasing Q_0 we start to lose control of the interaction intensity of a parton with a given x and $k_\perp \sim Q_0$ (and thus may err in the overall production rate). However, such partons do not branch any further, do not produce any soft offspring, so that the *shape* of the resulting energy distribution remains undamaged. Colour coherence plays here a crucial rôle.

It is important to realize that knowing the spectrum of *partons*, even knowing it to be a CIS quantity in certain sense, does not guarantee on its own the predictability of the *hadron* spectrum. It is easy to imagine a world in which each quark and gluon with energy k produced at the small-distance stage of the process would have dragged behind its personal "string" giving birth to $\ln k$ hadrons in the final state (the Feynman plateau). The hadron yield then would be given by a convolution of the parton distribution with a logarithmic energy distribution of hadrons from the parton fragmentation.

If it were the case, each parton would have contributed to the yield of non-

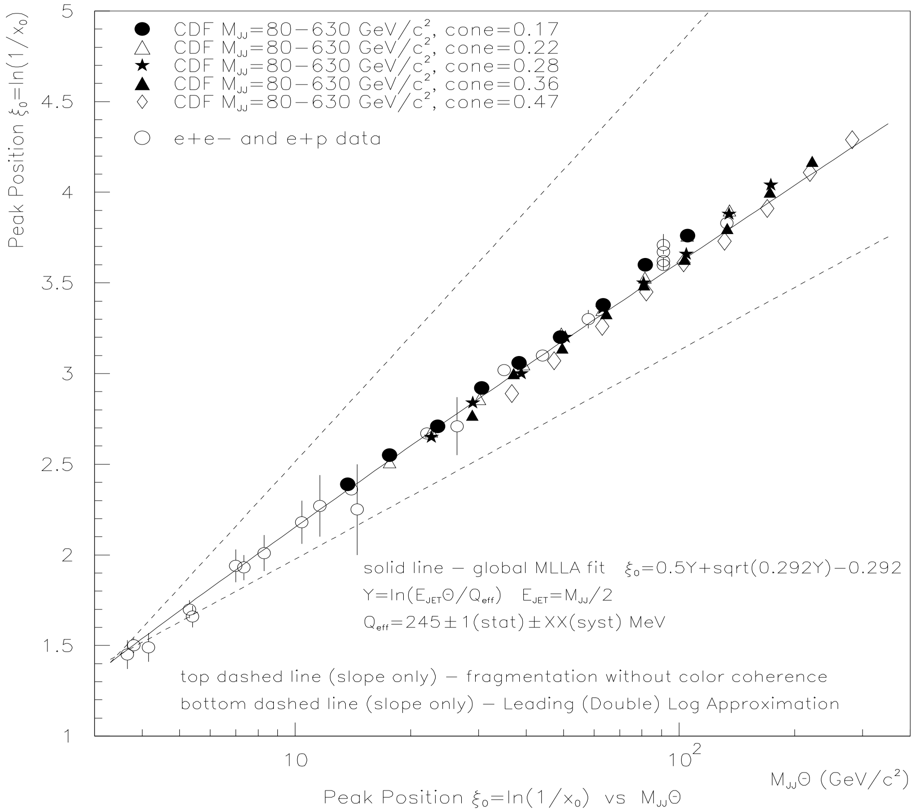

Figure 5. The position of the maximum versus the analytic MLLA prediction (Safonov 1999).

relativistic hadrons and the hadron spectra would peak at much smaller energies, $\xi_{\max} \simeq \ln Q$, in a spectacular difference with experiment.

Physically, it could be possible if the non-perturbative (NP) hadronisation physics did not respect the basic rule of the perturbative dynamics, namely, that of colour coherence.

There is nothing wrong with the idea of convoluting time-like parton production in jets with the inclusive NP parton→hadron fragmentation function, the procedure which is similar to convoluting space-like parton cascades with the NP initial parton distributions in a target proton to describe DIS structure functions.

What the nature is telling us, however, is that this NP fragmentation has a finite multiplicity and is *local* in the momentum space. Similar to its PT counterpart, the NP dynamics has a short memory: the NP conversion of partons into hadrons occurs locally in the configuration space.

In spite of a known similarity between the space- and time-like parton evolution pictures ($x \sim 1$), there is an essential difference between *small–x* physics of DIS structure functions and the jet fragmentation. In the case of the space-like evolution, in the limit of small Bjorken–x the problem becomes essentially non-perturbative and PT QCD loses control of the DIS cross sections (Mueller 1997, Camici &

Ciafaloni 1997). On the contrary, studying small Feynman–x particles originating from the time-like evolution of jets offers a gift and a puzzle: all the richness of the confinement dynamics reduces to a mere overall normalisation constant.

The fact that even a legitimate finite smearing due to hadronisation effects does not look mandatory makes one think of a deep duality between the hadron and quark-gluon languages applied to such a global characteristic of multihadron production as an inclusive energy spectrum.

Put together, the ideas behind the brave gluon counting are known as the hypothesis of Local Parton-Hadron Duality. Experimental evidence in favour of LPHD is mounting, and so is list of challenging questions to be answered by the future quantitative theory of colour confinement.

(d) QCD Radiophysics

Another class of multihadron production phenomena speaking in favour of LPHD is the so-called inter-jet physics. It deals with particle flows in the angular regions between jets in various multi-jet configurations. These particles do not belong to any particular jet, and their production, at the PT QCD level, is governed by coherent soft gluon radiation off the multi-jet system as a whole. Due to QCD coherence, these particle flows are insensitive to internal structure of underlying jets. The only thing that matters is the colour topology of the primary system of hard partons and their kinematics.

The ratios of particle flows in different inter-jet valleys are given by parameter-free PT QCD predictions and reveal the so-called "string" or "drag" effects. For a given kinematical jet configuration such ratios depend only on the number of colours (N_c).

For example, the ratio of the multiplicity flow between a quark (antiquark) and a gluon to that in the $q\bar{q}$ valley in symmetric ("Mercedes") three-jet $q\bar{q}g$ e^+e^- annihilation events is predicted to be

$$\frac{dN_{qg}^{(q\bar{q}g)}}{dN_{q\bar{q}}^{(q\bar{q}g)}} \simeq \frac{5N_c^2-1}{2N_c^2-4} = \frac{22}{7}. \tag{2.3}$$

Comparison of the denominator with the density of radiation in the $q\bar{q}$ valley in $q\bar{q}\gamma$ events with a gluon jet replaced by an energetic photon results in

$$\frac{dN_{q\bar{q}}^{(q\bar{q}\gamma)}}{dN_{q\bar{q}}^{(q\bar{q}g)}} \simeq \frac{2(N_c^2-1)}{N_c^2-2} = \frac{16}{7}. \tag{2.4}$$

Emitting an energetic gluon off the initial quark pair depletes accompanying radiation in the backward direction: colour is *dragged* out of the $q\bar{q}$ valley. This destructive interference effect is so strong that the resulting multiplicity flow falls below that in the least favourable direction transversal to the three-jet event plane:

$$\frac{dN_{\perp}^{(q\bar{q}\gamma)}}{dN_{q\bar{q}}^{(q\bar{q}g)}} \simeq \frac{N_C+2C_F}{2(4C_F-N_c)} = \frac{17}{14}. \tag{2.5}$$

At the level of the PT accompanying gluon radiation (QCD radiophysics) such predictions are quite simple and straightforward to derive. The strange thing is, that these and many similar numbers are being seen experimentally. The inter-jet particles flows we are discussing are dominated, at present energies, by pions with typical momenta in the 100–300 MeV range! The fact that even such soft junk follows the PT QCD rules is truly amazing.

(*e*) *Soft confinement*

Honestly speaking, it makes little sense to treat few-hundred-MeV gluons as PT quanta. What hadron energy spectra and string/drag phenomena are trying to tell us is that the production of hadrons is driven by the strength of the underlying colour fields generated by the system of energetic partons produced in a hard interaction. Pushing PT description down into the soft gluon domain is a mere tool for quantifying the strengths of the colour field.

Mathematical similarity between the parton and hadron energy and angular distributions means that confinement is very soft and gentle. As far as the global characteristics of final states are concerned, there is no sign of strong forces at the hadronisation stage which forces would re-shuffle particle momenta when the transformation from coloured quarks and gluons to blanched hadrons occurs. (For a recent review of MLLA-LPHD issues see Khoze & Ochs 1997.)

This observation goes along with what we have learnt from studying DIS, with special thanks to HERA which taught us that proton is truly fragile. It suffices to kick it with 1 GeV momentum transfer, or even less, and it is blown to pieces.

It seems that what keeps a proton together is not any strong forces between the quarks but merely quantum mechanics: the proton just happened to be the ground state with a given well conserved quantum number (baryon charge). It is interesting to see how easy is it to break a proton. To achieve that it is not even necessary to kick it hard. A soft scratch (or rather two) is enough to do the job.

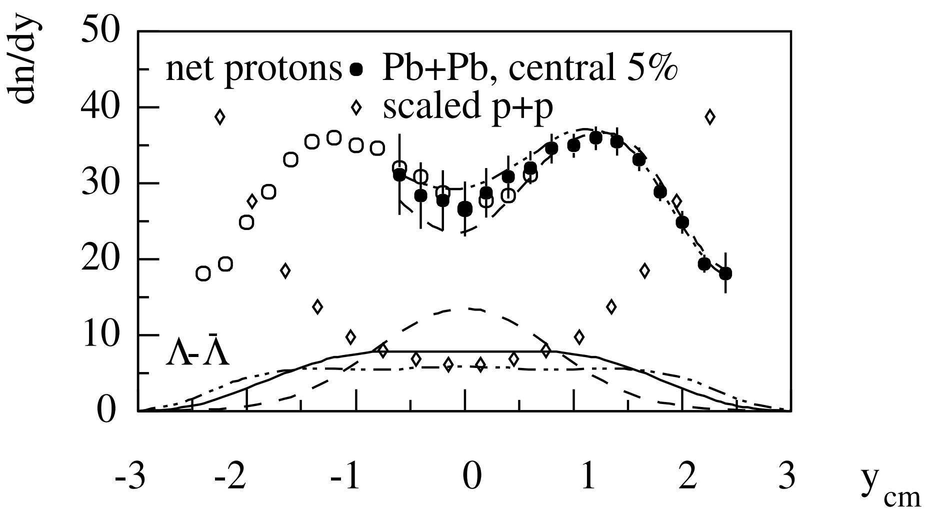

Figure 6. Proton “stopping” as seen at CERN by NA–49 (1999)

There is no sign of advocated fragility in a normal (minimum bias, soft) high energy proton-proton scattering. The famous leading particle effect shows that a projectile protons stays intact in the final state and carries away a major fraction of the incident momentum (the net proton spectrum "scaled $p + p$ " in Fig. 6). This is not surprising. In a typical pp interaction it is only one of the valence quarks of the proton that scatters. Internal coherence of the spectator quark pair remains undisturbed. In these circumstances the proton splits into a triplet quark and a spectator diquark which is in a colour anti-triplet state. On the hadronisation stage, the former picks up an antiquark and turns into a meson carrying, roughly, $z \simeq 1/3$ of the initial proton momentum, while the diquark (colour equivalent of a $\bar{q}$) picks up a quark forming a leading baryon ($z \simeq 2/3$).

It suffices, however, to organize a *double* scattering within a life-time of the intrinsic proton fluctuation in order to destroy the proton coherence completely (including that of the diquark which remained intact after the first scratch). Now the three quark-splinters of the proton separate as independent triplet charges and normally convert into three "leading" mesons (carrying $z \simeq 1/3$ each) in the final state. The proton decays, for example, into

$$p(1) \to \pi^+(1/3) + \pi^-(1/3) + K^+(1/3) + \ldots$$

with the baryon quantum number sinking in the sea.

This is what seems to be going on in the ion-ion scattering as shown in Fig. 6. Disappearance of leading protons is known as "stopping" in the literature. This I believe is an inadequate name: there is no way to *stop* an energetic particle, especially in soft interaction(s). Relativistic quantum field theory is more tolerant to changing particle identity than to allowing a large transfer of energy-momentum (recall relativistic Compton where the *backward* scattering dominates: electron turns into a forward photon, and vice versa).

If this heretic explanation of the "stopping" as proton instability is correct, the same phenomenon should be seen in the proton hemisphere of proton-nucleon collisions and even in pp. As we know, here there are leading protons. However, this is true on *average*. Even in pp collisions one can enforce multiple scattering (and thus full proton breakup) by selecting rear events, e.g. with larger than average final state multiplicity.

In all these cases (pp, pA, AB) "proton decay" should be accompanied by an enhanced strangeness production. Collecting experimental evidence in favour of proton instability is underway (Fischer 2000).

Soft hadronisation, likely absence of strong inter-parton forces, fragile hadrons — can it be reconciled with confinement in the first place? To the best of my knowledge, Gribov Super-Critical Light-Quark Confinement theory (GSCC) is the only scenario to offer a natural explanation to the puzzling phenomenology of multi-hadron production discussed above.

Light quarks are crucial for GSCC. If I was not so ignorant in theology (and was not brought up to hate philosophy), we could spend some time discussing, has it not been done on purpose that God supplied us with very light (practically massless u and d) quarks in order to make the hadron world easier to understand?

It is clear without going into much mathematics that the presence of light quarks is sufficient for preventing the colour forces from growing real big: dragging away a heavy quark we soon find ourselves holding a blanched D-meson instead. The light quark vacuum is eager to screen any separating colour charges.

The question becomes quantitative: how strong is strong? How much of a tension does one need to break the vacuum and organize such a screening? In the early 90-s Gribov has shown that in a theory with a Coulomb-like interaction between light fermions it suffices to have the coupling exceeding the critical value,

$$\frac{\alpha_{\text{crit.}}}{\pi} \simeq 1 - \sqrt{\frac{2}{3}}\,, \tag{2.6}$$

to have super-critical binding, restructuring of the PT vacuum, chiral symmetry breaking and, likely, confinement (Gribov 1999 and references therein). The word *super-critical* refers to the known QED phenomenon of so-called super-critical atoms. Dirac energy levels of an electron in a point-like Coulomb field of an ion with $Z > 137$ become complex. Classically, the electron "falls into the centre". Quantum-mechanically, it also falls, but into the Dirac sea: the ion becomes unstable and gets rid of an excessive electric charge by emitting a positron (Pomeranchuk & Smorodinsky 1945).

In the QCD context, with the colour factor $C_F = 4/3$ applied to the l.h.s. of (2.6), the critical coupling becomes

$$\left(\frac{\alpha_s}{\pi}\right)_{\text{crit.}} \simeq 0.137\,. \tag{2.7}$$

This number, apart from being easy to memorize, has another important quality: it is numerically small. Gribov's ideas, being understood and pursued, offer an intriguing possibly to address all the diversity and complexity of the hadron world from within the field theory with a reasonably small effective interaction strength (read: *perturbatively*).

3. Probing NP dynamics with PT tools

Can one talk about QCD coupling at small momentum scales? To answer such a question positively is not easy. Apart from courage, one needs to design some more or less definite prescription for quantifying an interaction strength at large distances where the very objects that are supposed to interact kind of don't exist! The best collection of arguments I could come up with, convincing or not, can be found in the proceedings of the HEP Vancouver conference (Dokshitzer 1998).

In recent years first steps have been made towards a joint technology for triggering and quantifying non-perturbative effects in CIS observables, both in "Euclid-translatable" cross sections and in the essentially Minkowskian characteristics of hadronic final states. The fact that the CIS observables are calculable in PT QCD (that is, remain finite when the collinear QCD cutoff μ is set to zero) does not imply that they are completely insensitive to NP dynamics. This only means that the genuine NP effects in CIS quantities manifest themselves as finite power suppressed corrections proportional to $(\mu^2/Q^2)^p \log^q(Q^2/\mu^2)$ with $p > 0$.

Simply by examining PT Feynman diagrams, one can find the exponents p, q for different observables. Knowing the leading power p is already useful: it tells us how (in)sensitive to confinement physics a given observable is.

More ambitious a programme aims at the *magnitudes* of power-suppressed contributions to hard cross sections and jet shape variables. (For references and a history of insights, phenomenological achievements and conceptual and numerical mistakes, this young subject is so rich with, the reader is invited to look into the proceedings of the Vancouver–1998 and Blois–1999 conferences.)

The magnitudes of the power-suppressed terms can be related with the behaviour of the coupling α_s in the infrared. Whatever the definition, it is thought to be a *universal* function that characterizes, in an effective way, the strength of the QCD interaction all the way down to small momentum scales. Given this universality, it becomes possible to *predict* the ratios of the Q^{-2p} contributions to observables belonging to the same class p.

In particular, the characteristic NP parameter

$$\alpha_0 = \frac{1}{\mu_I}\int_0^{\mu_I} dk\,\alpha_s(k^2), \quad (\mu_I = 2\ \mathrm{GeV}) \tag{3.1}$$

is conveniently used to quantify the NP hadronisation effects in CIS jet shapes, many of which belong to the $p=\frac{1}{2}$ class, i.e. exhibit large $1/Q$ power corrections. These include the thrust T, the so called C-parameter, invariant jet masses M_J and M_H (heavy-jet mass), the jet broadenings B_T and B_W (wide-jet broadening).

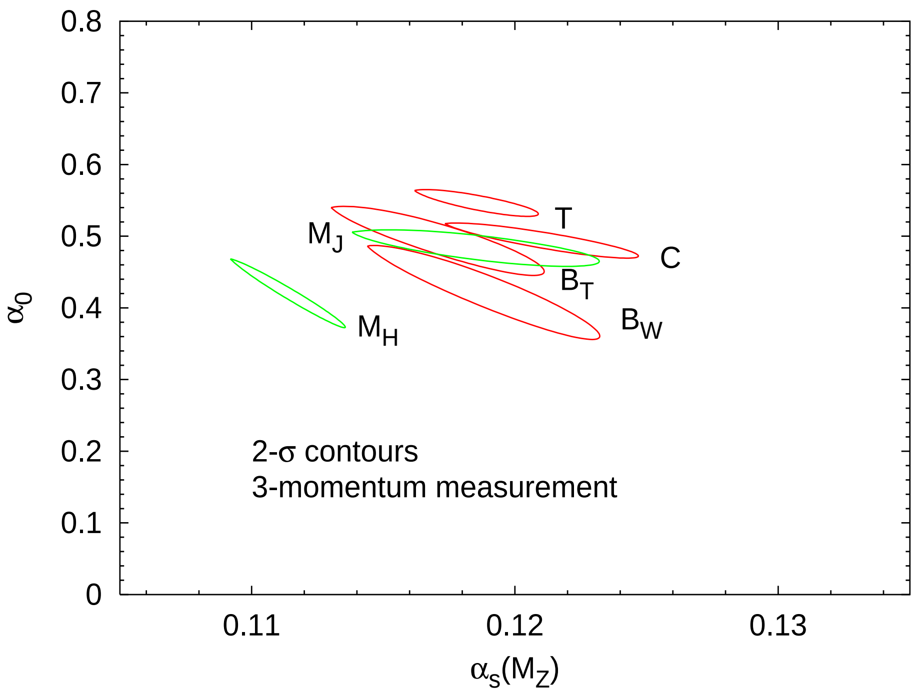

Figure 7. Infrared coupling from the means of jet shape variables (Salam & Wicke 2000).

Fig. 7 verifies the (in)consistency of independent experimental determinations of the PT and NP coupling from a variety of jet shapes, as it looks today (Salam & Wicke 2000). Given the relative weight of wishful thinking substituted for rigorous proofs in formulating theoretical rules of the game, you would agree that the hypothesis of universality (or, in other words, the notion of the universal infrared coupling) is not ruled out, to say the least. (NB: two of the displayed jet shapes, namely M_H and B_W, include jet selection, are therefore less inclusive and may have a reason to misbehave.)

Homework: Divide that α_0 by π ($\pi = 3$ would be good enough an approximation) and please compare with the Gribov critical coupling (2.7).

Shall we hear the bell ringing?

I am grateful to Gavin Salam for useful remarks. I want to congratulate Ian Butterworth, John Ellis and Erwin Gabathuler for the success of the Discussion Meeting they took a burden to have organized. I must seek their and the reader's forgiveness for a bad quality writeup. (Ruled by the celebrated Reference Theorem which states that *The quality of someone's paper is proportional to the number of references to your papers*, while *The quality of your own paper is inversely proportional to that.*)

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