Hard QCD

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Status of hard/perturbative QCD phenomena is briefly reviewed. Landau-Pomeranchuk-Migdal effect is discussed as a means for establishing links between particle and nuclear high-energy physics.

1. MESSAGES FROM THE HEP WORLD

Quantum Chromodynamics is the strangest of theories. On one hand, it *is* beyond any doubt the microscopic theory of the hadron world. Both the intrinsic beauty of QCD and the striking successes of QCD-based phenomenology speak for that. On the other hand, the depth of the conceptual problems that one faces in trying to formulate QCD as a respectable Quantum Field Theory (QFT) has no precedent in the history of modern physics. QCD nowadays has a split personality. It embodies "hard" and "soft" physics, both being hard subjects, and the softer the harder. (For more details on the present status of QCD, its problems and prospects, including nuclear issues, see [1].)

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 of massive lepton pairs, heavy quarks and their bound states, large transverse momentum jets in hadron-hadron collisions are classical examples of hard processes.

Perturbative QCD (pQCD) 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 pQCD approach, even to hard processes, is far from being obvious. One has to rely on plausible arguments (completeness, duality) and often to substitute Ideology for Theory.

Ideology is not necessarily a swear word (though my life-experience tends to tell me the opposite). An example of a good and powerful ideological concept is that of Infrared- and Collinear-Safety introduced by Sterman and Weinberg in the late 70's [2]. An observable is granted the status of infrared/collinear safety (ICS) if it can be calculated in terms of quarks and gluons treated as real particles (partons), without encountering either collinear

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($\theta \to 0$) 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). Given an ICS quantity, we expect its pQCD value predictable in the quark-gluon framework to be directly comparable with its measurable value in the hadronic world.

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 an ICS measure and stays under pQCD jurisdiction.

Speaking about the final state structure, we cannot predict, say, the kaon multiplicity of pion energy spectra. 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 pQCD — the jet pattern.

There are well elaborated procedures for counting jets (ICS jet finding algorithms) and for quantifying the internal structure of jets (ICS 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 visualises 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[1].

1.1. THERE ARE GLUON OUT THERE, AND THEY BEHAVE

LEP and SLAC e^+e^- experiments have reached a high level of sophistication. These days they study *identified* hadrons in *identified* (heavy quark-, light quark- or gluon-generated) jets by taking advantage of the prominent $Z^0 \to$ hadrons peak. Another QCD-factory is provided by the Fermilab $p\bar{p}$ Tevatron with its unique handle on jets with up to few-hundred-GeV transverse momenta. DESY HERA experiments scrutinise e^+p and e^-p DIS with an emphasis on small-x physics, which is bound to shed light on the transition region between hard and soft hadron phenomena.

pQCD does a spectacular job by covering many orders of magnitude in the basic jet production cross sections. Concerning the internal structure of jets, as well as multi-jet ensembles, the main message is that the perturbatively controlled secondary gluon radiation plays a crucial rôle. The bulk of final particle multiplicity is due to multiple radiation of "soft" gluons with relatively small momentum fractions $x \ll 1$. The structure of this radiation is determined by the geometry and colour topology of the underlying hard parton ensembles — the QCD antennae.

Gluons should produce soft gluon radiation more intensively than quarks, according to the celebrated ratio of "colour charges", $C_A/C_F = (N_c^2-1)/2N_c = 9/4$. As a result, the energy spectra of particles from gluon jets are expected to be softer, and the angular distributions broader as compared to quark-generated jets. These expectations are met by experiment. In Fig. 1 the angular profile of quark jets (DIS ZEUS, e^+e^- OPAL) is compared with that for gluon jets which dominate in $p\bar{p}$ (CDF, D0).

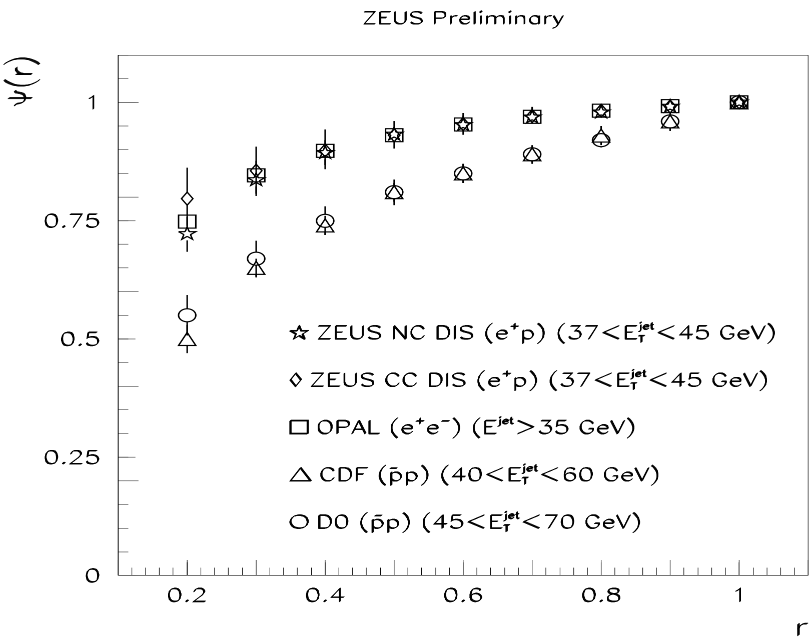

Figure 1. "Energy profile" of HERA quark jets. CDF/D0 (gluon) jets are broader.

Given the perfect identification of jets achieved by e^+e^- experiments, the C_A/C_F ratio was recently extracted from a comparative study of the scaling violation pattern in the fragmentation of quark and gluon jets. It can also be read out directly from the rate of growth of particle multiplicities with jet hardness, as shown in Fig. 2.

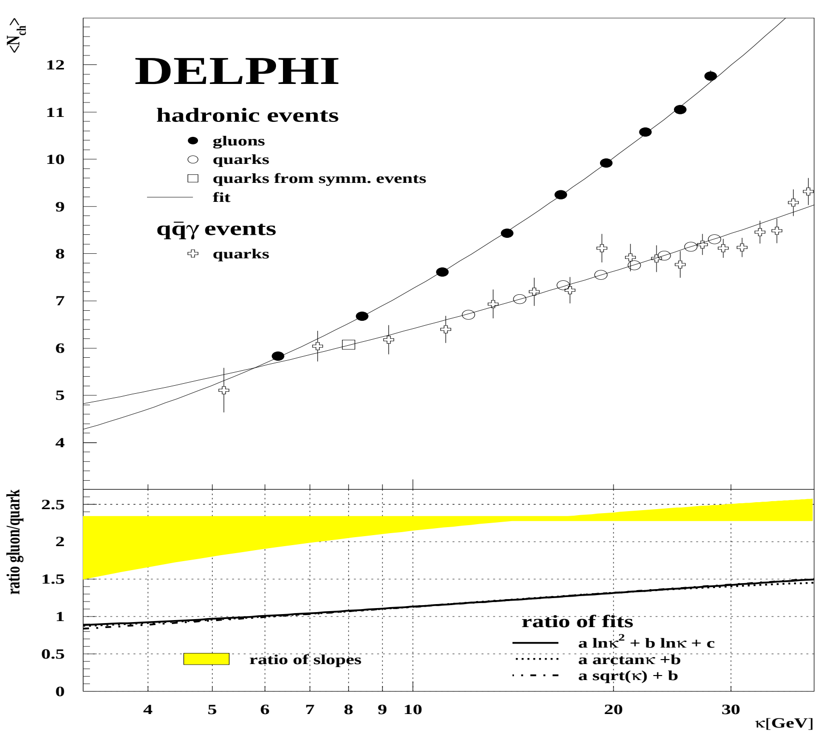

Figure 2. Charged hadron multiplicity from quark and gluon jets.

1.2. COHERENCE, LPHD AND SOFT CONFINEMENT

QCD coherence is essential in multi-gluon radiation. In order to formulate the parton multiplication processes in terms of probabilistic evolution, it is necessary to take into full account destructive interference effects. Coherence suppresses soft gluon emission at angles larger than the angular aperture of a bunch of hard radiating partons. With proper respect being payed to quantum-mechanical nature of radiation, the emerging cascade picture is based on the so-called angular ordering prescription for successive emissions of soft gluons [3].

Intra-jet coherence effects (angular ordering) are taken care of by smart MC event generators. At the level of analytic predictions, the corresponding technique is known as the modified leading logarithmic approximation (MLLA). It represents, in a certain sense, the resummed next-to-leading-order approximation. This step is necessary for deriving *asymptotically correct* predictions concerning multiple particle production in jets. This means that the MLLA parton-level predictions become exact in the $Q^2 \to \infty$ limit.

Gluon coherence inside jets leads to the so-called "hump-backed" plateau in one-particle inclusive energy spectra. The prediction was made (and the corresponding MLLA spectrum calculated with use of a 200-step Hewlett-Packard calculator) in 1983-84. Since then it has survived the LEP-1 scrutiny; more recently, it has been confirmed by the detailed CDF analysis; these days it is seen at HERA in the current fragmentation region formed by a struck quark in DIS.

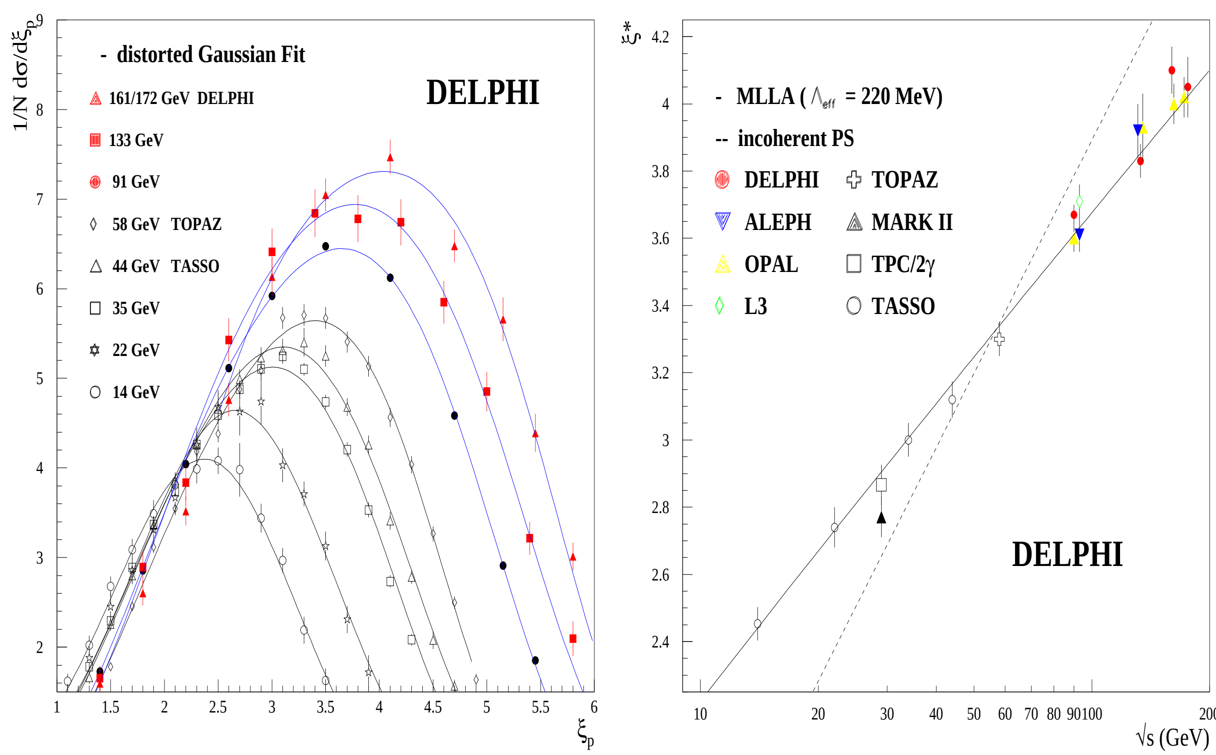

Figure 3. Hump-backed hadron spectrum, and predicted behaviour of its maximum.

Does the observation of the hump-backed plateau constitute a QCD test? Yes and no. On one hand, we do check the small-distance dynamics of coherent parton multiplication (this being in fact an expensive test of quantum mechanics). On the other hand, we gain important additional information about the non-perturbative dynamics of hadronisation: the similarity between the calculated parton and observed hadron spectra tells us that there is essentially no re-shuffling of momenta at the transformation stage from partons to hadrons. Such a property was envisaged and formulated as a hypothesis of local parton-hadron duality (LPHD, another example of an "ideological concept").

If the LPHD concept is correct, that is if hadronisation is *local* in configuration space, then the spectra of all hadron species should become asymptotically similar to each other and to that of partons (the bulk of which are relatively soft gluons). This leaves us with two global parameters to describe energy spectra of *hadrons* with $x_p = p/E_{jet} \ll 1$ at any sufficiently large $\sqrt{s} \equiv Q = 2E_{jet}$. They are: the scale of the perturbative QCD coupling $\Lambda_{\rm QCD}$ and the overall normalisation parameter K_g^h to recalculate the number of hadrons of a given species h from that of gluons. LPHD predicts, and experiments confirm, that the non-perturbative "conversion coefficients" K^h are true constants, depending neither on Q nor on the energy-momentum of the triggered hadron, x_p.

In Fig. 3 inclusive energy spectra of charged hadrons at different annihilation energies are shown ($\xi_p = \ln 1/x_p$).

What makes the story really surprising is that the pQCD spectrum is followed by pions (constituting 90% of charged hadrons in jets) with momenta well below 1 GeV!

A similar message comes from the study of coherent drag effects ("string effects") dealing with angular distributions of particle flows (multiplicity flows) *between* jets in hard events with non-trivial geometry/colour topology (*inter*-jet coherence). In particular, pQCD provides parameter-free predictions for the ratios of particle densities in different inter-jet valleys. At present energies these observables are dominated by pions in the 100–300 MeV momentum range. Amazingly, the distributions of this soft junk follow closely the coherent pQCD gluon radiation pattern. In other words, the pattern of the colour field that the underlying jet ensembles (hard parton antennae) develop.

We must conclude that hadronisation is a surprisingly soft phenomenon. As far as the global characteristics of final states are concerned, such as inclusive energy and angular distributions of particle flows, there is no visible change in particle momenta when the transformation from coloured quarks and gluons to blanched hadrons occurs.

A recent review of these and related topics can be found in [4].

2. LPM EFFECT: A BRIDGE BETWEEN HEP– AND HEN–PHYSICS

For many years QCD ideas have been used to picture high-energy scattering phenomena in nuclear matter. QCD-motivated constructions include the small-distance core of the intra-nuclear potential, modelling excitations of a nuclear target in terms of a colour tube, quark-gluon plasma, percolating strings, physics or chiral condensate, etc.

Historically, the nucleus has always been a primary source of inspiration for High Energy Particle (HEP) physics. Gribov's paper "Interaction of photons and electrons with nuclei at high energies" [5] laid a cornerstone for the concept of partons. Diffractive phenomena in hadron-nucleus scattering, and inelastic diffraction in particular, shed light on many a subtle problem of hadron interactions at high energies [6,7]. Multiple interactions in nuclear matter probe the internal structure of hadronic matter and make a nucleus serve as "colorometer" [8]. Thus, understanding High Energy Nuclear (HEN) phenomena is vitally important for developing QCD as a theory of the microscopic dynamics of hadrons.

In spite of this, rigorous applications of QCD to scattering in media are scarce, in the first place because of the complexity of the problems involved. The Landau-Pomeranchuk-Migdal effect is an example of such an application which addresses the issue of QCD processes in media "from the first principles" (if such a notion can be applied to QCD in its present state).

LPM is about radiation induced by multiple scattering of a projectile in a medium. In the QED context, Landau and Pomeranchuk noticed [9] that the energy spectrum of photons caused by multiple scattering of a relativistic charge in a medium is essentially different from the Bethe-Heitler radiation pattern. A few years later a quantitative analysis of the problem was carried out by Arkady Migdal [10]. Symbolically, the photon radiation intensity per unit length reads

$$\omega\frac{dI}{d\omega\, dz} \propto \frac{\alpha}{\lambda}\cdot\sqrt{\frac{\omega}{E^2}E_{LPM}}\;; \qquad \frac{\omega}{E} < \frac{E}{E_{LPM}}\,. \tag{1}$$

Here E is the energy of the projectile, and E_{LPM} is the energy parameter of the problem, built up of the quantities characterising the medium. They are the mean free path of the electron, λ, and a typical momentum transfer in a single scattering, μ (of the order of the inverse radius of the scattering potential):

$$E_{LPM} = \lambda\,\mu^2\,. \tag{2}$$

In QED the parameter E_{LPM} is in a ball-park of 10^4 GeV. Such an enormously large value explains why it took four decades to experimentally verify the LPM phenomenon [11].

The LPM spectrum should be compared with the standard Bethe-Heitler formula

$$\omega\frac{dI}{d\omega\, dz} \propto \frac{\alpha}{\lambda}\,, \tag{3}$$

which corresponds to independent photon emission for each successive scattering act.

Contrary to (3), the LPM spectrum (1) is free from an "infrared catastrophe": small photon frequencies are relatively suppressed, so that the energy distribution is proportional to $d\omega/\sqrt{\omega}$. Integrating (1) over photon energy ($\omega < E$ in the $E\to\infty$ limit), one deduces the radiative energy loss per unit length to be proportional to $\sqrt{E}$,

$$-\frac{dE}{dz} \propto \frac{\alpha}{\lambda}\sqrt{E\,E_{LPM}}\,. \tag{4}$$

In the QCD framework, expectations about energy losses due to gluon radiation off a *colour* charge propagating through a QCD medium were ranging, until recently, from constant, E^0, up to E^2.

The true answer is still $\sqrt{E}$ (for an "infinite" medium), as in the QED case, though the differential gluon energy spectrum proved to be very different from that of the LPM photons. In both problems the LPM phenomenon *suppresses* medium-induced radiation: a group of $N_{coh.}$ scattering centres, $N_{coh.} > 1$, acts as a single source of radiation.

In particular, the coherent LPM spectrum in QED can be presented as

$$\omega\frac{dI^{(LPM)}}{d\omega\, dz} = \omega\frac{dI^{(BH)}}{d\omega\, dz}\cdot\frac{1}{N_{coh.}^{(QED)}}\,, \qquad N_{coh.}^{(QED)} = \sqrt{\frac{E^2}{\omega\,E_{LPM}}} \,>\, 1\,. \tag{5}$$

To make a long story short, the QCD spectrum, amazingly enough, can be obtained from that for QED via a simple reciprocity relation [12], namely $\omega/E \Longrightarrow E/\omega$. This gives

$$N_{coh.}^{(QCD)} = \sqrt{\frac{\omega}{E_{LPM}}} > 1\,, \tag{6}$$

where E_{LPM} (2) is now of the order of 1 GeV. The gluon spectrum comes out to be E-independent and over-singular at small frequencies:

$$\omega\frac{dI}{d\omega\, dz} \propto \frac{\alpha_s}{\lambda}\cdot\frac{1}{N_{coh.}^{(QCD)}} = \frac{\alpha_s}{\lambda}\sqrt{\frac{E_{LPM}}{\omega}} = \alpha_s\sqrt{\frac{\mu^2}{\lambda}\cdot\omega^{-1}}\,; \qquad \omega > E_{LPM}\,. \tag{7}$$

At a semi-quantitative level, this result was obtained in [12] where the Gyulassy-Wang model for a QCD medium [13] was adopted. Further development [14] (BDMPS) included a treatment of finite-length media and the relation between the energy loss and jet broadening, apart from fixing the errors of the original treatment. The latter issue proved to be painful and slow a process. The final debugged set of the BDMPS predictions will soon become available, converging with that independently obtained by Bronislav Zakharov in the framework of an elegant functional integral technology [15].

Leaving the technical details aside, it is important to stress the main message that the study of the LPM phenomenon is sending us. Namely, that the physics of multiple interaction is infested with quantum mechanics, which makes the results often anti-intuitive and hardly accessible by means of classical probabilistic considerations.

First, and simplest of all, the fact that the formation time is finite plays a crucial rôle in the game. It is easy to accept that one should not treat multiple interaction of secondaries with the target as *independent* before a certain time $t = t_{form.}$ elapses. What is more difficult to digest, is that the very value of $t_{form.}$ depends on this interaction which cannot be modelled classically.

Imagine a relativistic quark traversing a QCD medium. The characteristic coherent length for induced gluon radiation, $\ell = \lambda \cdot N_{coh.}$, can be obtained from the following simple consideration. On one hand, the formation time of the radiation is

$$t_{form.} = \frac{\omega}{k_\perp^2}\,, \tag{8}$$

which is nothing but the time it takes to leave a "wave zone", in the language of the classical radiation theory. On the other hand, large formation times exceeding the mean free path, $t_{form.} \gg \lambda$, are essential for the LPM effect. This implies a random walk in the gluon transverse momentum, from one scattering centre to another,

$$k_\perp^2 \simeq \mu^2 \cdot N_{coh.} = \mu^2 \cdot \frac{t_{form.}}{\lambda}\,, \qquad t_{form.} = \frac{\lambda k_\perp^2}{\mu^2}\,. \tag{9}$$

Equating these two expressions for $t_{form.}$ we arrive at the correct expression (6) for $N_{coh.}$:

$$k_\perp^2 = \sqrt{\frac{\omega\,\mu^2}{\lambda}}\,, \qquad t_{form.} = \sqrt{\frac{\omega\,\lambda}{\mu^2}}\,; \qquad N_{coh.} = \sqrt{\frac{\omega}{\lambda\,\mu^2}} \equiv \sqrt{\frac{\omega}{E_{LPM}}}\; 1\,. \tag{10}$$

The transverse separation between a primary quark and a secondary gluon remains small as compared with the size of the scattering potential, $\Delta\rho_\perp \ll \mu^{-1}$ for $t < t_{form.}$. Therefore the quark-gluon system interacts with the medium as a whole, with the *quark* scattering cross section. Nevertheless, we have to treat the gluon transverse momentum as being accumulated in a course of successive independent scatterings of the gluon in the medium. To confuse you still more, the whole LPM spectrum emerges as a result of *interference* between the amplitudes of gluon emission at different times, say, at $t=t_0$ and $t=t_0+t_{form.}$. This means that during the formation time we have a quark-gluon pair in the amplitude, and a bachelor quark in the conjugate amplitude; so, is there a gluon or there is no? It is not a gluon but rather "to-be-a-gluon" that we are dealing with.

Another anti-intuitive prediction emerges when one takes into consideration the finite size of the medium [14]. Moving from the BH to the LPM spectrum we sliced a "brick" into groups consisting of $N_{coh.}$ scattering centres, each group acting coherently as an effective single centre. This implies that the longitudinal size of the medium, L, is large enough to embody at least one such a group:

$$L > t_{form.} = \lambda \cdot N_{coh.} = \sqrt{\frac{\omega\,\lambda}{\mu^2}}\,; \qquad \omega < \frac{\mu^2}{\lambda} L^2\,. \tag{11}$$

For a long medium, $L > L_{cr}$ with

$$L_{cr} \equiv \sqrt{E\,\frac{\lambda}{\mu^2}}\,, \tag{12}$$

this does not affect the energy loss. However, for a finite medium, $L < L_{cr}$, the situation changes. The largest medium-induced gluon energy becomes L-dependent, and we obtain the energy loss *per unit length* to be proportional to the *size of the medium*, instead of $\sqrt{E}$. Integrating over z leads to the total loss growing as L^2, a purely coherent enhancement effect which is difficult to digest "classically".

The ultimate trick that quantum mechanics plays with us in the LPM problem can be looked upon as an unexpected gift. One starts from an entirely "soft" environment: momentum transfer in a single scattering is small, and the very applicability of the perturbative consideration is far from secure. In spite of this, the result proves to stay under perturbative control: it is multiple scattering which ensures the dominance of small distances, and thus the applicability of pQCD. What matters in the problem is the *accumulated* transverse momentum, $k_\perp^2 = \mu^2 \cdot N_{coh.} \gg \mu^2$, which stays large provided $\omega \gg E_{LPM} \sim 1$ GeV.

To see how this actually happens, let us write down the full answer for the differential radiative energy loss ($L < L_{cr}$):

$$-\frac{dE}{dz} = \frac{\alpha_s(B^2)\,N_c}{8}\,\frac{\mu^2\,\tilde{v}(B^2)}{\lambda} \cdot L\,. \tag{13}$$

Here B is the impact parameter inverse proportional to the accumulated $k_\perp$,

$$B^2 \simeq \frac{\lambda}{\mu^2} \cdot L^{-1} \;\ll\; (1\,\mathrm{fm})^2\,. \tag{14}$$

The dimensionless factor $\tilde{v}$ characterises the scattering potential:

$$\tilde{v}(B^2) = \frac{1}{\sigma\,\mu^2} \int_0^{1/B^2} dQ^2\, Q^2\, \frac{d\sigma}{dQ^2}\,, \tag{15}$$

where $d\sigma/dQ^2$ is the differential single scattering cross section for a given momentum transfer Q. We notice that all three entries, i.e. the mean free path λ, the radius of the potential μ^{-1} and the total scattering cross section σ in $\tilde{v}$ are ill-defined quantities, since they are determined by soft physics (properties of finite-momentum-transfer scattering in a medium). However, they enter in (13) (and into the differential gluon energy distribution) in a combination which is dominated by small distances. Indeed, invoking the definition of the mean free path, $\lambda^{-1}=\rho\sigma$, with ρ the density of centres, we arrive at

$$\frac{\mu^2\,\tilde{v}(B^2)}{\lambda} \;=\; \rho \int_0^{1/B^2} dQ^2\, Q^2\, \frac{d\sigma}{dQ^2}\,. \tag{16}$$

In QCD the integral on the right-hand-side is logarithmically enhanced, being dominated by large momentum transfers,

$$\frac{d\sigma}{dQ^2} \propto \frac{\alpha_s^2}{Q^4}\,, \qquad \text{for} \quad \mu^2 \;\ll\; Q^2 \;\ll\; B^{-2} = \mu^2\,\frac{L}{\lambda}\,, \tag{17}$$

and therefore stays under pQCD control (at least in the logarithmic approximation).

Finally, let us mention an interesting relation between the radiative energy loss and the "broadening" of the transverse momentum distribution of the projectile due to scattering in a medium. The *width* of the transverse momentum distribution, $p_{\perp W}^2$, is proportional to L (random walk) and is determined by the same parameter (16). Therefore the following relation between the density of energy loss and the jet broadening holds, which is noticeably independent of the interaction dynamics and of the colour of the projectile[14]:

$$-\frac{dE}{dz} \;=\; \frac{\alpha_s\,N_c}{8}\,p_{\perp W}^2\,. \tag{18}$$

This result is in accord with a general inequality derived earlier by Brodsky and Hoyer[16].

Induced radiation should affect propagation of partons both in the initial and in the final states of hard interactions in a medium.

At present there seems to be a problem of reconciliation of experimental findings: the A-dependent transverse momentum received by *incoming* partons [17] appears to be much smaller than that received by *outgoing* ones [18]. Broadening of Drell-Yan lepton pairs (IS effect) points at a value of the parameter (16), for cold nuclear matter, 10–15 times *smaller* than the corresponding value characterising FS effects. The latter was extracted by Luo, Qiu and Sterman [19] from the analysis of the $p_\perp$-disbalance of two jets produced in $\gamma A \to 2$ jets at Fermilab. One thing seems however certain: in a *hot* QGP the expected broadening and energy loss are still bigger. In particular, a QGP with temperature $T=250$ MeV should produce LPM gluons with energies and transverse momenta up to

$$\omega_{\max} \;\simeq\; 250\,\text{GeV}\left(\frac{L}{10\text{ fm}}\right)^2, \qquad \left(k_\perp^2\right)_{\max} \simeq p_{\perp W}^2 \;\simeq\; 5\,\text{GeV}^2\,\frac{L}{10\text{ fm}}\,, \tag{19}$$

while the corresponding values for cold nuclear matter are estimated to be *at least* twice (more realistically, by a factor 20–30) smaller.

LPM physics supplies QCD jets produced inside, and propagating through, a medium, with extra gluons with a quite narrow, and weird, angular distribution $\Theta \sim (\omega_0/\omega)^{3/4}$ ($\omega_0 \simeq 500$ MeV for the hot matter). It also forces initial partons to lose energy prior to engaging into a hard interaction (to produce a Drell-Yan pair, large-$p_\perp$ jets, etc.). The former effect makes the jets softer, broader and more populated. The latter should cause a medium-dependent "factorisation breaking" by driving the hard cross sections away from the A^1 regime. Medium-induced scaling violation effects will be more pronounced near the phase space boundaries where a relatively small energy loss matters.

Study of these issues has begun. We can expect the present-day discrepancy to be clarified in a foreseeable future, and more reliable estimates of the QCD LPM effects to become available, in particular, concerning the comparison of cold and hot media.

3. FEEDBACK: HEN $\Longrightarrow$ HEP

There is no doubt that the HEN physics should, and eventually will, teach its HEP counterpart. We cannot expect HEN physics to be able to clarify many a smoking-gun issue, which is necessary to combat our ignorance about the hadron dynamics. However, some specific HEN phenomena should provide indispensable tools for digging out crucial information about the structure of hadrons and their interactions, inaccessible otherwise. The only problem is, how to locate such specific phenomena. To this end, an obvious strategy would be to concentrate on unexpected/unexplained things happening.

To name a few, an excess of small-mass lepton pairs [20], Hagedorn-type particle abundances [21], baryon stopping, large Λ/p and, especially, $\bar{\Lambda}/\bar{p}$ ratios [22], not to mention jumpy J/ψ nuclear absorption [23,24].

3.1. SOME MODERATELY NASTY REMARKS

As an ignorant outsider, I am allowed a couple of heretic comments concerning the ways some of the above-mentioned puzzles are being discussed.

To start with, the LPM physics sends a warning message: a classical Monte-Carlo modelling of intra-nuclear particle multiplication at very high energies is like robbing banks — tempting but dangerous.

Two more comments are of purely linguistic nature. First, I would restrain from using the word "temperature" in the discussion of the famous exponential fit to hadron abundances, $N \propto \exp(-m/T)$, or $\exp(-m_\perp/T)$, if that matters. The relative hadron yields from $e^+e^- \to$ hadrons satisfy the same phenomenological law, $p : K : \pi \simeq 1 : 2 : 20$. However, the e^+e^- final states are anything but "thermal": produced secondaries are far apart in configuration space and get no chance to talk to each other. At least in this clear case, the Hagedorn "temperature" is a universal property of the *vacuum*, of the parton$\to$hadron transition, rather than a thermal characteristic of a particle ensemble.

Finally, the sacred "deconfinement" itself. To avoid confusion we'd rather call it an "ultimate confinement". Indeed, if we define confinement as a problem of preventing massless colour fields from appearing in the physical spectrum of the theory, than the plasma phase does the job. Massless gluons disappear due to complete colour screening. In other words, gluons are "confined" simply because they have no chance to propagate freely, being scattered off the colour charges in a medium. Photons acquire mass in metals because they are screened ("confined", if you wish) not because they are "deconfined".

3.2. LOOKING INTO PUZZLES

From the study of diffractive pA phenomena we have learnt that the proton projectile can be caught in a "squeezed", "transparent" state which penetrates a nucleus (for more details, see [1]). Its counterpart — a swollen proton — can be characterised as a configuration with larger than typical relative distances between the quarks. In such a state "strings" are stretched, colour (or, pion) fields are stronger, the vacuum is virtually broken. This is the proton-"perpetrator" which, contrary to the proton-"penetrator", should willingly interact with the target.

Along these lines, a nice physical explanation of baryon stopping and hyperon production phenomena was suggested by Kharzeev [25]. "Baryon junction" is the name of the game [26], however we can do without. To understand the physics of the matter it suffices to keep in mind that 1) each constituent quark sits in a colour-triplet state, and 2) high-energy scattering is a coherent phenomenon.

The proton as a QFT object is a coherent sum of various configurations. The quantum portrait of a projectile, its field fluctuation, stays frozen in the course of interaction at high energies, so that each fluctuation scatters independently [6]. A minimum-bias pp scattering is normally a peripheral process, in which one of the constituent quarks in the proton is hit. Coherence between a struck quark and the rest of the projectile gets broken, and the system decays. Roughly 2/3 of the initial proton momentum goes into the forward baryon, the remaining 1/3 being shared, successively, by secondary mesons and baryons (in the ratio $\approx 8:1$) which form a hadron plateau. The latter can be viewed as breaking of a standard $3\otimes\bar{3}$ "colour string". It is important to stress that the coherence between the two spectator quarks is undisturbed, so that a peculiar $3\otimes 3\otimes 3$ colour structure of the proton remains hidden (the diquark acting coherently as $3\otimes 3=\bar{3}$).

Eventually, with a probability of about 1/8, the first breakup of the string produces a sea diquark–anti-diquark $d\bar{d}$ instead of a $q\bar{q}$ pair. In such a case a baryon emerges that contains the struck quark with 1/3 momentum fraction, $B_{1/3}=q_{1/3}q_0q_0$. The rest of the system has a $q_{1/3}q_{1/3}\bar{q}_0\bar{q}_0$ content and *predominantly* decays into two mesons, $M_{1/3}+M_{1/3}$. The $B_{2/3}+\bar{B}_0$ option is formally open as well. However, this channel should be strongly suppressed, the reason being the statistics of the energy levels which a *coherent* quantum mechanical system chooses to occupy (recall the "Hagedorn exponent").

Turning to a *nuclear* target we get a chance to successively scratch *two* constituents in a swollen proton configuration, provided the energy is large enough so that the two scatterings occur within a Lorentz-dilated life-time of the fluctuation. By doing so we fully break the coherence of the initial system. In such circumstances we have to consider breakup of three independent $3\otimes\bar{3}$ strings attached to a common "junction point" somewhere inside the proton. With probability $\sim (7/8)^3 \simeq 67\%$ such a system will decay into three leading mesons, each carrying roughly a third of the proton momentum, $M_{1/3}+M_{1/3}+M_{1/3}$ and a "stopped" baryon, $+B_0$. Moreover, a simple quark-pickup picture enriches the Λ/p ratio for the latter: combinatorially, the probability to have *at least* one strange quark in the "centre" can be estimated as $1-(2/3)^3\simeq 70\%$.

This simple picture produces the most spectacular prediction for the central yield of anti-baryons. By allowing one string to break up "baryonically", that is via $d\bar{d}$, we shall have $B_{1/3}+M_{1/3}+M_{1/3}$ as leading hadrons. An adjacent central quark soup consisting of $q_0q_0\bar{q}_0\bar{q}_0$ is no longer coherent, contrary to the single-scratch pp case discussed above.

Therefore, there is no reason to expect a $B_0 + \bar{B}_0$ channel to be exponentially smaller than the meson one. (Moreover, the colour structure of the soup, $q \otimes q = \bar{3}$, $\bar{q} \otimes \bar{q} = 3$, cries for the creation of a baryon pair!) Leaving aside an open question of the absolute yield of central anti-baryons, let us concentrate on the statistics of strangeness.

Three vacuum-produced quarks (suppressing the prefix anti-) form 27 flavour combinations falling into $8 + 12 + 6 + 1$ states with strangeness 0,1,2 and 3, respectively. The $SU(3)_{flavour}$ nomenclature of these states is

# of states	decuplet	octet	singlet	final state	$\to n$	$\to p$	$\to \Lambda$
8=	$\Delta^{++}\Delta^{+}\Delta^{0}\Delta^{-}$	$+2\cdot(p\,n)$			2+2	2+2	0+0
12=	$\tilde{\Sigma}^{+}\,\tilde{\Sigma}^{0}\,\tilde{\Sigma}^{-}$	$+2\cdot(\Sigma^{+}\,\Sigma^{0}\,\Sigma^{-})$	$+3\cdot\Lambda$		0+3	0+1	3+5
6=	$\tilde{\Xi}^{0}\,\tilde{\Xi}^{-}$	$+2\cdot(\Xi^{0}\,\Xi^{-})$			0+0	0+0	2+4
1=	Ω^{-}				0+0	0+0	1+0

The right half of the table describes the approximate n, p, Λ composition of the final state after (strong, weak, radiative) decays. The first components in the three rightmost columns separate the yields from the decuplet states Δ, $\tilde{\Sigma}(1385)$, $\tilde{\Xi}(1530)$ and Ω.

These statistics apply both to the Λ/p and $\bar{\Lambda}/\bar{p}$ ratios. In the former case the *observed* value will be smaller because of stopped initial protons feeding the denominator. The anti-baryon ratios come out clean,

$$\bar{\Lambda}/\bar{p} \simeq \frac{6+9}{2+3} = 3\,, \qquad \bar{n}/\bar{p} \simeq \frac{2+5}{2+3} = 1.4\,.$$

(Note that the first ratio does not depend on the relative weight of the decuplet states.)

4. INSTEAD OF CONCLUSION

QCD at present is still in rather limited, tough no longer a primitive, stage of development. It is bound to use the language of quarks and gluons, that is to talk perturbatively, and is trying to extend its grip, from the solid base of hard processes, to the realm of soft hadron interactions. On an A–B plot, pQCD is steadily gaining grounds in the origin, $A \cdot B = 1$, and tries to crawl along the $A \cdot B = A$ line. How about a big jump into the $A \sim B \sim 10^2$ spot?

The physics of ion-ion high energy collisions being so abundantly rich, the only hope of a big success lies within a clever strategy for extracting a needle from a hay-stack. What one desperately needs here is a constructive way of moving from a b-strategy (sit upon, and try to feel) to the m-strategy (make use of a magnet).

One obvious magnet is QGP searches (whether you believe this particular needle being in there or not). The QGP state is usually thought to be formed in the collision, which provides a melting pot for individual nucleons. A large E_t-yield, for example, is considered to be a sign of the phase transition into such a state, in the course of the collision.

There is however an alternative way of looking upon things. Studying the *tails* of various distributions, small cross sections, we start probing rare configurations of colliding objects. Thus, very large E_t may be looked upon as a *precondition* for the collision, rather than the result of it: to observe a larger than typical transverse energy yield, we catch the colliding nuclei *pre-prepared* in a rare, confinement-perpetrating, virtually melted configuration à la the desired plasma state. In such configurations the yield of lepton pairs should be higher (extra pions, or antiquarks, around); a reduced yield of J/ψ (heavy traffic) is to be expected, strangeness may start to "misbehave", etc.

The word *coherence* has been haunting us through the text. Quantum mechanics cannot be *understood*. The best one can do is to get used to it. The HEP community rediscovered quantum mechanics, in the QCD context, at the beginning of 80's [3]. It is never too late.

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REFERENCES

1. Yu.L. Dokshitzer, *EPS HEP Conference*, Jerusalem, August 1997; hep-ph/9801372.
2. G. Sterman and S. Weinberg, Phys.Rev.Lett. 39 (1977) 1436.
3. for review see Yu.L. Dokshitzer, V.A. Khoze, A.H. Mueller and S.I. Troyan, *Basics of Perturbative QCD*, ed. J. Tran Thanh Van, Editions Frontières, Gif-sur-Yvette, 1991.
4. V.A. Khoze and W. Ochs, Int.J.Mod.Phys. A12 (1997) 2949.
5. V.N. Gribov, ZhETF 57 (1969) 1306 [Sov.Phys.JETP 30 (1970) 709].
6. E.L. Feinberg and I.Y. Pomeranchuk, Suppl.Nuov.Cim. III (1956) 652.
7. H. Miettinen and J. Pumplin, Phys.Rev.Lett. 42 (1979) 204.
8. L.L. Frankfurt, G.A. Miller and M. Strikman, Ann.Rev.Nucl.Part.Sci. 44 (1994) 501.
9. L.D. Landau and I.Ya. Pomeranchuk, Dokl. Akad. Nauk SSSR 92 (1953) 535, 735.
10. A.B. Migdal, Phys.Rev. 103 (1956) 1811.
11. P.L. Anthony et al., Phys.Rev.Lett. 75 (1995) 1949.
12. R. Baier et al., Phys.Lett. 345 (1995) 277.
13. M. Gyulassy and X.-N. Wang, Nucl.Phys. B 420 (1994) 583.
14. R. Baier et al., Nucl.Phys. B 483 (1997) 291; ibid. B484 (1997) 265.
15. B. Zakharov, JETP Lett. 63 (1996) 952; ibid. 65 (1997) 615.
16. S.J. Brodsky and P. Hoyer, Phys.Lett. B 298 (1993) 165.
17. D.M. Alde, et al., Phys.Rev.Lett. 66 (1991) 2285.
18. T. Fields and M.C. Corcoran, Phys.Rev.Lett. 70 (1993) 143;
    R.C. Moore, et al., Phys.Lett. B 244 (1990) 347;
    M.D. Corcoran, et al., Phys.Lett. B 259 (1991) 209.
19. M. Luo, J. Qiu and G. Sterman, Phys.Rev. D 49 (1994) 583.
20. J.Wambach, *in these Proceedings.*
21. F.Becattini, *in these Proceedings.*
22. P. Braun-Munzinger, *in these Proceedings.*
23. L. Ramello, *in these Proceedings*;
    M. Gonin, *Proc. of the Quark Matter '96 Conf., Eds. P. Braun–Munzinger et al.*, Nucl.Phys. A 610 (1996) 404.
24. D. Kharzeev, *in these Proceedings.*
25. D. Kharzeev, Phys.Lett. B 378 (1996) 238; S.F. Vance, *in these Proceedings.*
26. L. Montanet, G.C. Rossi and G. Veneziano, Phys.Rep. 63 (1980) 149.