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P 3a Dynamics of Hadronic Reactions

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§1. Introduction

The subject of this talk is a very wide one: I am supposed to cover the theoretical developments which have recently occurred in our understanding of hadronic processes, *i. e.*, of processes controlled by strong interactions.

Such reactions include, besides conventional low p_T hadron physics and large p_T hadron-hadron collisions, also processes with initial and/or final leptons, such as $e^+e^- \to$hadrons, $l+p \to l+X$ and $pp \to l^+l^- + X$.

In this talk I shall not be able to describe, nor even mention, all the interesting contributions in the field. Instead, and, I believe, in agreement with the guide-lines of the Program Committee, I shall try to give an over-all view of the present "state of the art," emphasizing here and there a few new results by which I have been particularly impressed.

It is perhaps the first time, at this type of Conference, that "soft" and "hard" hadron physics are discussed together in a plenary talk. I took this as a suggestion to try to present both aspects of hadron physics as one body of knowledge, stressing, as much as possible, connections, analogies and differences.

The plan of the talk will be as follows: After giving a panoramic overview of "hadronland," and of the talk, I shall discuss, in order, hard processes, low-energy (spectroscopy), but very incidentally, intermediate to high energies at low p_T, and superasymptotic energies at low p_T. For each of these subjects, I shall refer to other speakers of parallel or plenary sessions for more details and/or complementary information. I shall conclude with a brief, tentative comparison of soft and hard hadron physics.

I will address myself, in particular, to some of the basic questions in hadron physics that have occupied our minds over the past few years. These are:

1) Is it conceivable that Quantum Chromodynamics (QCD) is at the basis of hadron physics? Can we actually test QCD?

2) Does QCD conflict a priori with other successful (though incomplete) descriptions of strong interactions, such as the Regge–Mueller approach, the dual string model, Reggeon Field Theory, Regge-bootstrap schemes, or does it rather provide a unifying link among them and with the constituent (parton) models of hard processes?

3) What do hard and soft processes have in common?

Concerning the question of what is the *present* evidence for QCD, I shall rather refer you to a recent paper of Bjorken,[1] concluding that such evidence is far from established.

Certainly, the most convincing way to answer those questions would be to prove (or disprove!) confinement in QCD and then to compute the hadronic spectrum from few input parameters. This would be certainly convincing, but, in spite of some nice progress recently made,[2] it could still take a little too long.

Meanwhile, an alternative, less satisfactory, but cheaper way to answer those questions can be the following:

a) Extract from QCD as many as possible testable "predictions" (*i. e.*, results believed to be unaffected by the confinement mechanism) and check them against available data. Find "clean" tests of QCD.

b) Look at the general structure of QCD with the aim of relating it to those other (partially) successful approaches to soft hadronic phenomena.

c) See if and how soft and hard processes are related in QCD, and in nature.

Incidentally, if this less ambitious approach should disprove QCD, then we could spare ourselves from the harder task of proving confinement.

To anticipate the conclusions, what seems to come out of such an analysis is the following:

i) With a little faith in the gentle behaviour of the confinement mechanism, a lot of predictions can be made for a variety of hard processes, some absolute, some relating different reactions. So far, QCD looks all right, but certainly not proven[1] (see also the following talk by Field). Some stringent tests are near, such as the p_T^{-4} behaviour of $d\sigma/dp_T^2$ predicted by QCD (with its normalization relative to deep inelastic data) which should soon show up,[3] and the peculiar predictions of QCD for e^+e^- produced hadronic jets in the PETRA–PEP energy region.[4] The need for clean tests is certainly felt.

ii) The topological graph structure of QCD assuming, of course, confinement to hold, suggests a clear relation to dual, Gribov and even to Regge bootstrap theories (*e.g.*, the strong coupling is fixed by QCD). It looks to be just a matter of finding, in each regime, the relevant collective degrees of freedom. This may force us into QCD-inspired semi-phenomenological theories, at least for the near future.

iii) QCD points at some crucial differences between hard and soft hadronic phenomena. Progress is underway towards understanding some intermediate regimes and this should help in finding the connection, if any. It looks that a common denominator for hard and soft hadron physics may exist, as I shall explain at the end of this talk.

§2. A Panoramic Overview of Hadronland

We have schematically represented in Fig. 1 the various regimes of hadron physics on a two-dimensional map with energy and momentum transfer (actually their logarithms) giving the co-ordinates.

First of all there is an unphysical region ($2p_T > \sqrt{s}$, but it actually extends in the complex planes of these variables). It is in the asymptotic part of this region (the deep Euclidean region) that improved perturbation theory (IPT) in the running coupling constant $\alpha_s(Q^2)$ can be justified for asymptotically free theories, such as QCD.

More interesting to us is, of course, the deep physical (Minkowski) region where hard processes take place. We shall discuss in a moment the use of IPT in this region, noticing that the use of IPT is confined at present to finite (and not too small) values of $x = 2p_T/\sqrt{s}$ ($Q^2/2M\nu$), hence to a strip along the diagonal. Ordinary low p_T physics also lies inside a strip, this time parallel to the energy axis with the width of such a strip shrinking if the diffraction peak does so. Asymptotically in this second strip, is where Reggeon Field Theory (RFT) is a popular theoretical framework; coming down, we encounter other interesting descriptions such as the concept of a bare Pomeron pole (ISR-Fermilab regions?), that of an exchange degenerate Reggeon and, finally, the resonance description of hadronic reactions. We shall see later how these descriptions can be linked to one another and this will also explain various arrows and words in the picture.

The resonant region is common to both strips I have mentioned. What happens there we believe to be related to what happens at larger values of E and/or p_T either by the old Dolen–Horn–Schmit or by the Bloom-Gilman duality relations.

The most prominent (and sad) feature of Fig. 1, however, I find to be the fact that so much of our hadronland lies *outside* these two strips (about which we think we have some understanding). This is the region that should provide the link between hard and soft hadron physics, but little is known about it. There are interesting speculations, due primarily to Feynman, that valence partons are related to Reggeon exchange and sea partons and gluons to Pomeron (vacuum) exchange; but as we go to higher E and p_T ("semi-hard" processes) we are quite stuck in the dark. Nevertheless, something has been moving on in this direction during last year, and I shall try to report on it.

Before moving on, let me say that the scale of E and p_T relevant to hadrons cannot be predicted from QCD: it is a free parameter given to us by nature: it is, say, 500 MeV. Until 1974 we thought this to be the only relevant scale, but the discovery of new heavy quarks has put us into some sort of puzzle.

Turning now to Fig. 2, I am anticipating there what seems to emerge from QCD for the effective degrees of freedom relevant to the various regimes depicted in Fig. 1.

We go from a parton-like description of hard processes with quarks and gluons, colour and flavour as explicit degrees of freedom to a dual-string-type representation of the resonance

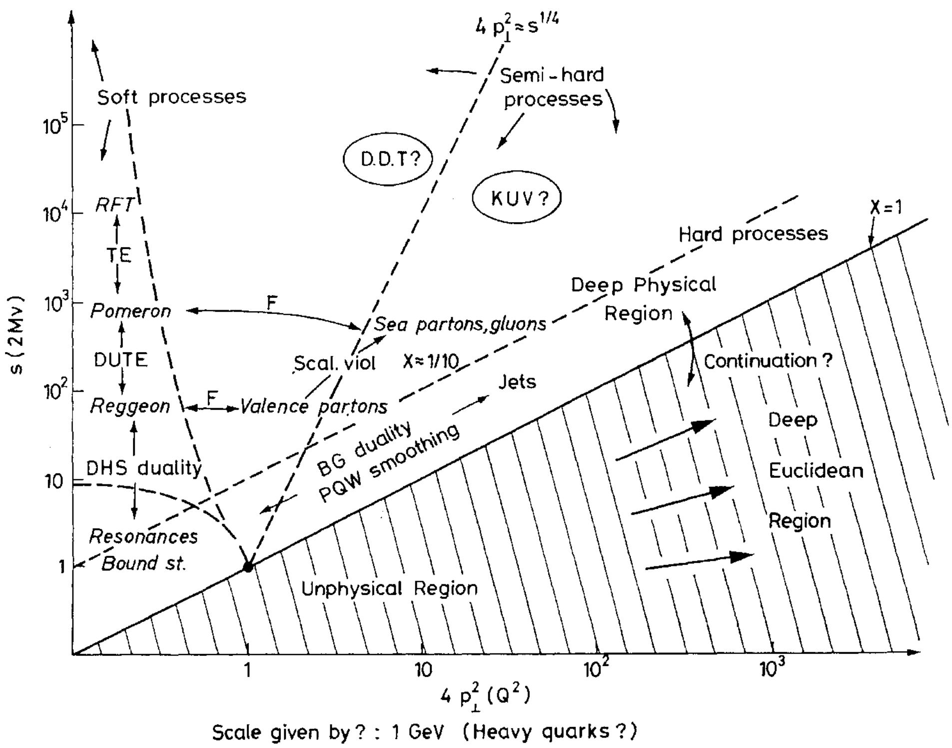


Fig. 1. An overview of "hadronland" and of its present understanding.

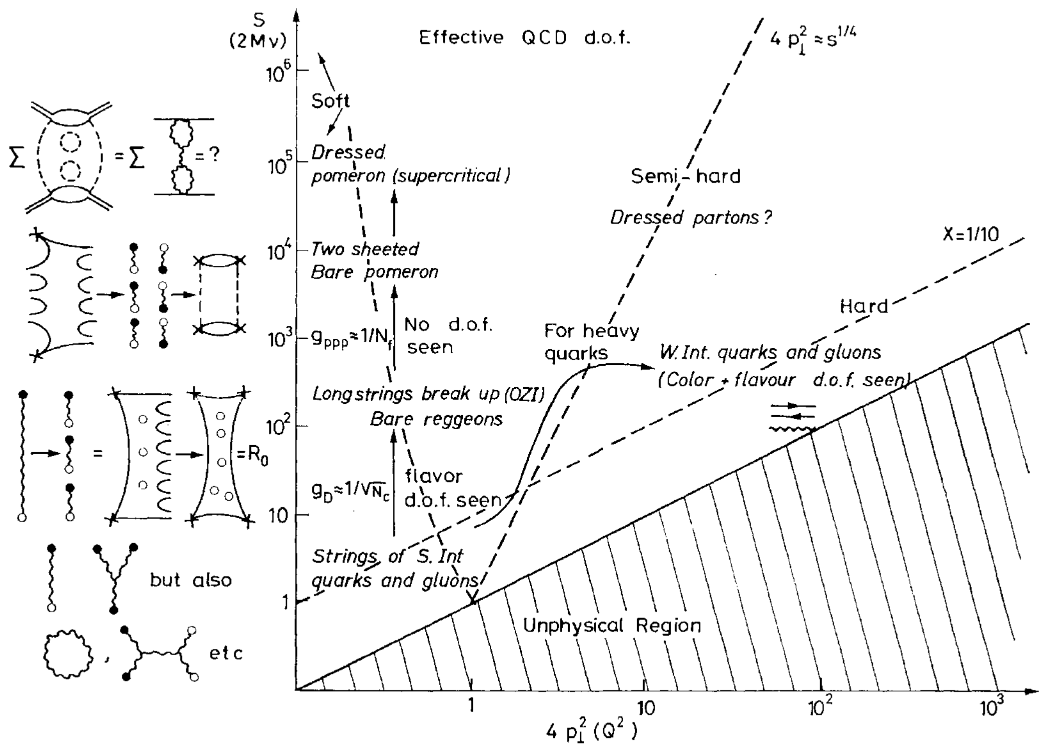


Fig. 2. Degrees of freedom appropriate to the various regions of "hadronland."

region in which colour ceases to be an explicit degree of freedom, since hadrons are supposedly colourless. Yet the existence of the underlying (hidden) colour degrees of freedom can be shown[5,6] to control the stability of our states ($g_{\rm strong} \simeq 1/\sqrt{N_c}$ if N_c is the number of colours). As we increase the energy, the excited, long strings break and, through multiperipheral dynamics, lead to a new degree of freedom, the exchange degenerate, ideally mixed planar Reggeon. Then, as we increase the energy further, the flavour singlet component takes over quantum number exchange and, by the time we are at Fermilab-ISR energies, or even before, the new effective degree of freedom will be vacuum exchange, *i. e.*, the bare Pomeron $\boldsymbol{P}_0$. Flavour degrees of freedom have also been lost at this stage, but there is a trace of them in the couplings of $\boldsymbol{P}_0$ (*e.g.*, $g_{P_0P_0P_0}\sim 1/N_f$ if N_f is the number of flavours[6]). In the end, the tube-shaped bare Pomerons interact with each other to give higher topological structures, an example of which is shown again in Fig. 2. Summing these corresponds to solving Reggeon Field Theory (RFT), presumably in the so-called supercritical region.

Finally we can ask what are the degrees of freedom in the "semi-hard" region. The authors of ref. 7 have suggested that they are still quarks and gluons, but no longer as structureless point-like objects: partons will get dressed and their predictable form factor will be measured by experiments in these kinematical regimes.

We are now ready for going down into hadronland and visit some of its most interesting sites.

§3. Hard Processes Involving Hadrons

3.1. *Strictly hard processes*

Under this name we mean to include both hadron and lepton initiated reactions as long as they are hard, *i.e.*, all invariants are large and of the same order. It does not look theoretically meaningful, indeed, to treat separately the reactions $l^+l^-\rightarrow$hadrons, $l+h\rightarrow l+$hadrons, $h+h\rightarrow$hadrons, $h+h\rightarrow l^+l^-+$hadrons (l for lepton, h for hadron) as long as they are all hard.

It is not an exaggeration to say that there has been a lot of progress on this subject, not only since Tbilisi, but even more within last year. Here I shall concentrate on the essential theoretical points and refer you to Politzer[8] for more details and to Close and Field[9] for the phenomenological applications, actual numbers, checks with data and so on.

In short, the trend has been towards an increased confidence in the use of improved perturbation theory (IPT) for asymptotically free (AF) theories, and in particular for QCD, outside the range of light-cone-dominated processes [*i.e.*, $\sigma_T(e^+e^-\rightarrow$hadrons) and $\sigma_{\rm incl}(l+h\rightarrow l+$hadrons)]. The way to go about it had been to deal with infra-red (IR) insensitive quantities, *i. e.*, quantities which, perturbatively, have a smooth, finite limit as the gluon and/or the quark mass goes to zero. This is because the renormalization group (RG) equation, which is always valid, plus AF relates a large Q^2 problem to a small coupling ($\alpha_s(Q^2)\sim\log^{-1}Q^2/\Lambda^2$), small mass ($m(Q^2)\sim m\cdot\Lambda/Q$) problem. In general, the usefulness of a small $\alpha_s(Q^2)$ is upset by large IR logarithms of Q^2/m^2. If those are absent (IR insensitivity) an asymptotic expansion in $\alpha_s(Q^2)$ may be all right, in the same sense as it is used in QED.

I will now sketch a general approach to this type of question, which relies heavily on the classic work[10] of Kinoshita and of Lee and Nauenberg (KLN). Those old results, originally obtained in QED, look to be valid for QCD or for any other renormalizable (but not super-renormalizable) field theory. They have, indeed, a very simple physical meaning which I shall now try to convey.

Consider a cross-section $\sigma(i\rightarrow f)$ which is finite in lowest order. Then, although higher order corrections make $\sigma(i\rightarrow f)$ IR divergent, the sum

$$\sum_{i'\sim i}\sum_{f'\sim f}\sigma(i'\rightarrow f')$$

where $i'(f')$ are degenerate with $i(f)$, is completely free of divergences. There are two types of divergences to cancel and correspondingly two types of degenerate states to be considered, *i. e.*:

a) *"Soft"* divergences, due to $m_{\rm gluon}$ or $m_{\rm photon}\rightarrow 0$ and to the possibility of emission of soft quanta by massive charged states. These are, for instance, the only divergences occurring in massive (*i.e.*, $m_e\neq 0$) QED. In QCD they are cured by adding cross-sections to *degenerate final* states such as q, $q+$soft gluon, $q+2$ soft gluons, etc.

b) *"Collinear"* divergences, due to the decay of a massless particle (a gluon or a massless quark) into two collinear, but hard, massless particles. According to KLN these divergences, which will be our main concern here, are cured by adding cross-sections with *initial* and final degenerate states consisting of a collection of hard massless collinear quanta (which is also a massless system).

The reason why divergences of type (a) will not bother us is that any interesting hard cross-section will automatically sum over soft bremsstrahlung processes.

At this point, a simple classification of hard QCD processes* follows from the KLN theorem. It looks as follows[11]:

1) *Processes with no initial coloured quanta* (*e.g.*, $e^+e^- \to$hadrons).

1a) If we do not detect individual hadrons in the final state, and instead limit ourselves to total or jet-inclusive cross-sections (these latter as defined by Sterman and Weinberg[12]), then KLN says that such processes are free of IR problems. Hence, as emphasized in ref. 12, cross-sections for producing n hadronic jets in e^+e^- collisions are completely calculable in QCD (absolute normalization included) and exhibit scaling violations only through their expansion in powers of $\alpha_s(Q^2)$.

For further developments along these lines, see ref. 4, where a whole set of possible QCD tests at PETRA–PEP is presented.

1b) If we look instead at a single particle spectrum $\sigma(e^+e^- \to h(p)+X)$, fixing the momentum p, or better $x=2|p|/\sqrt{Q^2}$, will automatically include soft bremsstrahlung (which does not change x) but will *not* include hard collinear bremsstrahlung (which brings you down in x). Hence this process is *not* IR finite by KLN, is not absolutely calculable, and, as one can easily show, will exhibit "large"** scaling violations.

2) *Processes with one initial coloured quantum* (*e.g.*, deep inelastic scattering).

2a) If we are totally inclusive, or jet inclusive à la Sterman and Weinberg,[12] there are no IR problems associated with the final states. There are, however, divergences due to the initial quark state which is *not* accompanied by the full set of degenerate states demanded for cancellation by KLN. As a consequence (and as we know in this case from operator product expansion methods), the process is not absolutely calculable and exhibits "large" scaling violations (anomalous dimensions).

2b) If, on top, we are also detecting a final hadron with well defined x, new IR divergences and scaling violations occur, just as in case 1(b).

3) *Two incoming coloured quanta* (*e.g.*, $q\bar{q} \to \gamma^*+X$, $qq \to qq$, $q+g \to q+g$, etc.)

3a) If we are jet-inclusive we now pick up IR divergences from both initial quanta. Since these are related to some extra hard collinear quarks and gluons in the initial state and since the two original incoming quanta are not themselves collinear, it is not surprising that we get (see below) separate IR singular factors for each incoming quantum and consequently *factorized* scaling violations.

3b) New scaling violations occur as in 1b) and 2b) if final particle inclusive spectra are looked at.

What is all this good for? Well, it looks as if it can provide, for the first time, a convincing derivation of the much used and successful parton model, or better of that particular version of it which incorporates QCD-predictable scaling violations.

This nice result came out of a large number of papers which appeared this year on the subject, but one should not forget that some of the basic points had been already laid down by the pioneering works of Gribov and Lipatov[13] and of Mueller.[14]

The new interest in the subject has been triggered by the paper of Sterman and Weinberg[12] and by those of Politzer.[15] Technically, the recognition that, by the KLN theorem,[10] the problem can be reduced to the study of collinear IR divergences and that those are easiest to study in physical gauges (in particular in the axial gauge, which is free of unphysically polarized gluons and of ghosts, and which has the Ward Identity relation $Z_1=Z_2$ as in QED), have been crucial developments.

Progress has been fast: from the one-loop calculation of some processes[12,15,16] to that of a general process[11]; from leading log calculations at higher orders in a few processes,[1,17] to their extension to arbitrary reactions[18] and, finally, to the analysis (if not the explicit calculation) of all non-leading log s.[18,19]

The result of all that has been to show that:

a) Soft divergences indeed cancel, reducing the singularities from double log s for each power of α_s to single log s.

* All these processes ought to be considered at the elementary constituent level. It is believed that, if a soft hadronic wave function can be defined, the results will apply to the actual world with only trivial modifications (see example below).

** Here, and in the following, by "large" scaling violations we just mean $(\log Q^2)^\gamma$ dependences typical of an anomalous dimension γ, as opposed to the asymptotically vanishing violations encountered in $\sigma(e^+e^- \to$hadrons).

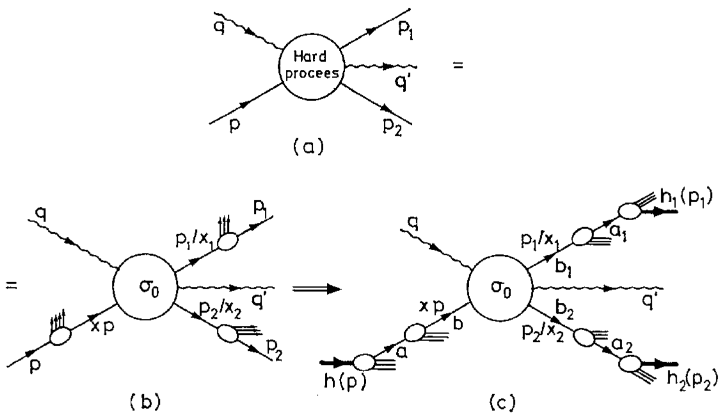


Fig. 3. QCD analysis of a typical hard process (γ^*+hadron→hadron$_1$+hadron$_2$+γ^*+X) in the leading log Q^2 approximation.

b) Collinear divergences are universal, *i.e.*, process independent, and factorize, *i.e.*, IR dependence factors from Q^2 dependence as originally conjectured by Politzer[15]. Furthermore there is an independent IR divergent factor for each incoming and detected final coloured quantum.

c) The above IR divergences can be absorbed into universal Q^2 independent factors to be identified with structure functions and fragmentation functions at some reference point Q_0^2.

d) Q^2 dependent factors (scaling violations) are also universal and predictable by QCD.

e) Different processes can be compared by IPT because suitable ratios of hard cross-sections are IR insensitive.

f) Last, but not least, the resulting diagrammatic understanding of OPE results seems to open the way to further extensions (see below).

In order to show that the results (a)-(f) correspond to our present ideas about the QCD parton model, let me illustrate the result in an example at the leading log level.

Consider the process of Fig. 3a to be translated eventually into a hadronic process such as that of Fig. 3c, where q, q' are electromagnetic or weak currents. One finds, in accord with KLN, that large Q^2 dependent factors (scaling violations) come separately from each coloured initial state and detected final state (Fig. 3b, at the parton level where the small blobs stand for collinear bremsstrahlung processes). When hadrons are added in, the final expression corresponding to Fig. 3c reads:

$$\begin{aligned}&\sigma(h(p)+\gamma^*(q)\to h_1(p_1)+h_2(p_2)+\gamma^*(q')+\mathrm{x})\\&\quad=\sum_{b,b_1,b_2}\int_0^1 \mathrm{d}x\mathrm{d}x_1\mathrm{d}x_2 G_h^b(x,Q^2)\\&\qquad\times D_{b_1}^{h_1}(x_1,Q^2)D_{b_2}^{h_2}(x_2,Q^2)\times \qquad (1)\\&\times\sigma_0(b(xp)+\gamma^*(q)\to b_1(p_1/x_1)+b_2(p_2/x_2)+\gamma^*(q'))\end{aligned}$$

where σ_0 is the lowest order hard cross-section and the structure (fragmentation) functions $G(D)$ have a well-defined Q^2 dependence, the one predicted by QCD, and are process independent (as long as the process is hard). This is exactly the QCD parton model prescription.

I will conclude this point by mentioning that the extension to non-leading log s is an important step forward, first because one is never too sure about how much trust can be put in leading log calculations and, second, because, at present values of Q^2, p_T, etc., such non-leading terms appear to be important.[20]

The lesson here seems to be that, although for each process the computation of non-leading log s is highly complicated,[21] when relating different hard processes much simpler results are obtained.[18]

Finally, I should point out that Mueller[22] has recently given an OPE-like justification of the above diagrammatic results adding thus further confidence in their validity.

3.2. *Semi-hard processes: DDT[7] extension*

The considerations made in 3.1 hold for inclusive processes in which all invariants are large and their ratios are all of order one. We now want to consider, following ref. 7, an intermediate regime which we call "semi-hard."

Consider, for instance, the (constituent level) reactions:

i) $q\bar{q}\to l^+l^-(q_T)+X$

ii) $\gamma^* q\to q(q_T)+X$

iii) $\gamma^*\to q(p_1)+\bar{q}(p_2)+X;\ (p_1-p_2)_T=q_T,$

where q_T is: in i), the transverse momentum of the lepton pair; in ii), the transverse momentum of the final quark relative to the direction of the "struck" quark (which can be determined); and finally, in iii), q and $\bar{q}$ lie on opposite side jets and q_T is the transverse momentum of the slower of the two relative to the direction of the faster.

The region we would like to consider is the one in which:

$$\Lambda^2 \ll q_T^2 \ll |Q^2| \tag{2}$$

where Q^2 is the virtual photon momentum and Λ the usual hadronic scale parameter ($\Lambda\simeq$ 500 MeV). In other words, we want both $\alpha_s(Q^2)$ and $\alpha_s(q_T^2)$ to be small but, unlike the case of real hard processes, we shall not be able to neglect terms of order $\alpha_s \log(Q^2/q_T^2)$. As I explained in § 2, this is a very large and interesting region of phase space.

Let us consider, to be definite, the Drell–Yan process i). If, instead of fixing q_T, we would evaluate $\langle q_T^2\rangle$, the result would be

$$\langle q_T^2\rangle \simeq Q^2\alpha_s(Q^2)\simeq Q^2/\log Q^2 \tag{3}$$

due to the fact that α_s is dimensionless. In terms of QCD, axial gauge diagrams, the process giving rise to this large $\langle q_T^2\rangle$ are those of Fig. 4a. Since the radiated hard gluons carry transverse momentum up to a finite fraction of Q, eq. (3) follows together with the fact that the annihilating $q\bar{q}$ pair can be anywhere from on shell to $O(Q^2)$ off shell.

Having instead restricted q_T to have a fixed, small value (relative to its average value) we are forcing those emitted gluons to have small transverse momenta and, consequently, we are fixing the off-shellness of the annihilating pair to be $O(q_T^2)$.

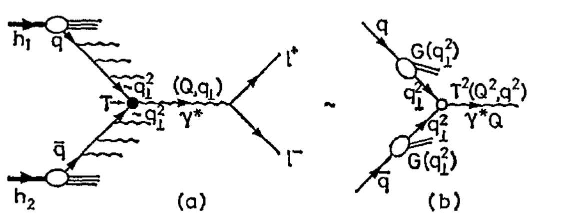


Fig. 4. Schematic understanding of the DDT formula.

DDT then claim to have obtained, within some more restricted region than (2), a very simple expression for the differential cross-section, which reads (DDT formula):

$$\frac{d\sigma}{dQ^2 dy dq_T^2}=\frac{4\pi\alpha^2}{9sQ^2}\sum_f e_f^2\frac{\partial}{\partial q_T^2}$$
$$\times[G_{h_1}^f(x_1, q_T^2)G_{h_2}^f(x_2, q_T^2)T^2(Q^2, q_T^2)]; \tag{4}$$
$$x_1\cdot x_2=Q^2/s;\ x_1/x_2=e^y,$$

where, essentially, the G structure functions, evaluated at q_T^2, can be understood as the result of the evolution of the quark density up to the q^2 value of the annihilating pair and the T factor is the *quark form factor* associated with the electromagnetic vertex (see Fig. 4b).

DDT have estimated this form factor and found that it is related (but not identical) to the Sudakov form factor.[23] They find

$$T\simeq\int_1^2 dz\exp(-B/2z^2)>\exp(-B/2)\equiv S, \tag{5}$$

where S is the Sudakov form factor:

$$B=-c_F\frac{\alpha_s(q_T^2)}{\pi}\log^2 Q^2/q_T^2;\ c_F=4/3. \tag{6}$$

Notice that T is a more "gentle" form factor than S.

A few remarks are in order about the DDT formula, eq. (4):

1) If naively integrated in q_T^2 from 0 to $O(Q^2)$ it clearly reproduces the Drell–Yan formula (with scaling violations, of course).

2) As a result of the form factor T, the q_T distribution of the lepton pair is flatter than the naïve expectation $(1/q_T^2)\alpha_s(q_T^2)$.

3) In a sense, one is measuring the form factor of the quark. Whereas in hard processes partons behave as point-like objects, in these semi-hard processes they behave as *dressed* particles.

4) The DDT derivation is not very simple and clear. It would seem important to have some double check of their simple final result.

5) Finally, one is dealing with a leading log approximation and some control of the non-leading terms would also be desirable.

3.3. Looking inside QCD jets

I would now like to describe briefly another extension of perturbative techniques for hard processes, made possible through our diagrammatic understanding of scaling violations.

This work[24] too goes in the direction of somewhat softer physics, because it deals with multiparton (and multiparticle) spectra inside the same (quark or gluon) jet. At first sight, it looks surprising that such quantitity as a two-parton spectrum inside a jet can be computed because one thinks that the invariant mass $(p_1+p_2)^2$ of the system ought to be small. This, however, is not the case, and again, as in eq. (3), one knows that*

$$\langle(p_1+p_2)^2\rangle \sim x_1x_2\cdot Q^2/\log Q^2. \quad (7)$$

The main results are the following:

1) In the leading log approximation, multiparton spectra inside a jet can be computed in terms of the tree diagrams of an effective (non-local) ϕ^3 theory. The resulting "jet calculus" rules are very simple.

2) This prescription satisfies many consistency checks (*e.g.*, momentum and charge conservation sum rules). In six-dimensional ϕ^3 theory (ϕ_6^3 is also AF) it gives, as a non-trivial by-product, a recent result of Taylor.[25]

3) For single particle QCD spectra the main results are:

i) The ratio of the average number of gluons (quarks) in a quark jet to the average number of gluons (quarks) in a gluon jet is finite and given by:

$$\frac{\langle n_i\rangle|_q}{\langle x_i\rangle|_g} \to c_F/c_A=4/9;\ i=q,\bar{q},g. \quad (8)$$

This result seems to prove a rather old conjecture of Brodski and Gunion.[26] Furthermore, the fact that the ratio does not depend on the type of "detected" parton makes us believe that the result simply extends to hadron multiplicities.

ii)

$$\frac{\langle x_i\rangle|_g}{\langle x_i\rangle|_q} \to c_F/c_A=4/9;\ i=q,\bar{q},g, \quad (9)$$

where $\langle x_q\rangle_q$ is the average x of a quark in a quark jet etc..

iii) The spectrum near $x=1$ is softer in a gluon jet than it is in a quark jet by a factor $(1-x)$, up to log $s(1-x)$. Properties i), ii) and iii) clearly support the general feeling[9] that gluon jets are softer and that they yield higher multiplicities.

A somewhat related, interesting observation, due to Shizuya and Tye,[27] is that the Sterman–Weinberg opening angle[12] for a gluon jet, δ_g, is much bigger than the corresponding one for quark jets, δ_q. One finds:

$$\delta_g=(\delta_q)^{c_F/c_A} \gg \delta_q;\ c_F/c_A=4/9. \quad (10)$$

All these results raise the question of whether gluon jets will be easy to see experimentally [see ref. 4) for an optimistic view on this point].

4) For two parton QCD spectra the main results are:

i) For x_1x_2 finite, various limits can be considered, *e.g.*,

$$(1-x_1-x_2)\ll 1;\ (1-x_1-x_2)\ll(1-x_1)\ll 1$$

etc.. Precise predictions, somewhat reminiscent of double Regge fragmentation behaviour emerge. They confirm the previous statements on gluon jets being softer than quark jets.

ii) For $x_1\sim 0$, x_2 finite, the two partons become essentially uncorrelated.

iii) For $x_1, x_2\sim 0$ there are long-range (in log x) correlations and one finds

$$\langle n(n-1)\rangle \simeq c\langle n\rangle^2, \quad (11)$$

where c depends on the type of jet considered, and also on the observed partons, and $c>1$. The constant c is also somewhat dependent on the way the IR singularities are cut off. On the other hand, all such dependence disappear when we compare quark and gluon jets. One gets, in particular

$$[\langle n(n-1)\rangle-\langle n\rangle^2]|_{q\,jet} = \frac{4}{9}[\langle n(n-1)\rangle-\langle n\rangle^2]|_{g\,jet} \quad (12)$$

independently of the observed species (hence valid for hadron?).

5) In general one finds, independently of the regularization procedures used,

$$p(z_i)|_{q\,jet}=c_F/c_A\,p(z_i)|_{g\,jet};\ c_F/c_A=4/9, \quad (13)$$

where $p(z)$ is the Feynman gas pressure as a function of the chemical potentials $z_i(i=q, \bar{q}, g)$. Furthermore, KNO scaling follows from the general structures of the theory.

6) No phenomenon à la Cornwall–Tiktopoulos[28] is seen to occur. One possibility is

* This is also why the results for parton spectra described below can be transformed into hadronic spectra using the only input of fragmentation functions (see ref. 24)).

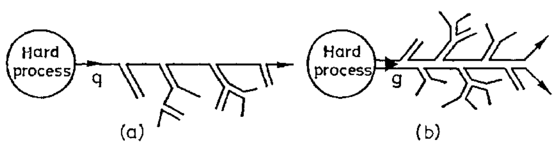


Fig. 5. Difference between a quark (a) and a gluon (b) jet in the large N limit. The graphical notation is that of refs. 5 and 6.

that things work out as in ϕ_6^3, where Taylor has found[25] for normalized exclusive cross-sections $\hat{\sigma}_n = \sigma_n/\sigma_T$:

$$\hat{\sigma}_n = \hat{\sigma}_0(1-\hat{\sigma}_0)^n;\ \hat{\sigma} \xrightarrow[Q^2/m^2\to\infty]{} 0;$$

$$\hat{\sigma}_n/\hat{\sigma}_0 \to 1;\ \sum_n \hat{\sigma}_n = 1. \qquad (14)$$

This can be recognized as an ideal Bose gas distribution (Planck spectrum) corresponding to a black body with

$$\frac{kT}{\hbar\omega} \simeq (\log Q^2/m^2)^{1/9}$$

as can be checked from the explicit behaviour[25] of $\hat{\sigma}_0$.

In a "naïve" regularization method for the QCD IR problem (which ϕ_6^3 does not have) one finds that the problem of computing the generating function of QCD jets is reduced to that of a so-called Markov branching process. This is exactly the problem encountered in studying the evolution in time of the population of competing species[29] (here quarks and gluons) having fixed probabilities of mutation and birth (but no death).

In this case one finds[24] that the gluon jet gives again a Planck spectrum (except that now $kT/\hbar\omega$ is IR divergent) and the q jet gives the spectrum of C_F/C_A independent black bodies. Two amusing limits of this result are worth mentioning. For $N_c \to \infty$ $C_F/C_A \to 1/2$ and the result can be understood by the fact that, in the large N limit, only planar diagrams survive[5,6] giving (see Fig. 5) a "one-sided" jet for the quark (5a) as opposed to a "two-sided" jet for the gluon case (5b). The other interesting case is the QED limit, corresponding to $C_F/C_A \to \infty$. In such a case the quark jet is infinitely many independent black bodies: this is well known to give a Poisson distribution, the standard QED result.

I should mention, however, that, in a different (dimensional type) regularization of the IR problem, which at present looks on more solid grounds, one gets, for instance:

$$D \equiv \sqrt{\langle n(n-1)\rangle - \langle n\rangle^2} = \gamma\langle n\rangle; \qquad (15)$$

$$\gamma = 1/\sqrt{3} \text{ for } g \text{ jet}, \quad \gamma = \sqrt{3}/2 \text{ for } q \text{ jet}$$

which disagrees with the ideal gas prediction (but agrees, as usual, in the ratio of quark to gluon jet).

Amazingly, the result eq. (15) for a gluon jet is exactly the known empirical Wrobleski relation[30] which is seen to hold in pp hadrons at low p_T. This could support models, like the one of Pokorski and Van Hove,[31] in which gluons are responsible for pionization in hadronic collisions. It could also support, however, topological expansion models of the Pomeron (see below). The quark jet prediction should be checked in e^+e^- collisions.

7) Finally, it has been possible to combine the techniques of refs. 7 and 24 in order to gain some more differential information (*e. g.*, q_T spread) on final particles inside QCD jets. Interesting results appear to come out, like again a broadening of the naïve q_T spectrum, but I have no time for going into more detail here.

3.4. *Planarity and duality in hard processes*

It turns out that when QCD diagrams are computed in physical (*e.g.*, axial) gauges, the leading diagrams for collinear IR divergences are the planar ones (see, *e.g.*, refs. 7 and 18). One can then argue that non-planar contributions to hard processes are down for two reasons:

i) Powers of $\alpha_s(Q^2)$, *i.e.*, inverse powers of $\log Q^2$ (this is only so in the axial gauge).

ii) Inverse powers of N_c or N_f, this being the case in any gauge.[5,6]

As a result planar diagrams in the axial gauge will completely dominate the dynamics of hard QCD processes.

It looks also natural to assume that, if confinement takes place at all in QCD, that should occur already at the planar level. As discussed below, this provides a very suggestive link between QCD and dual string theories. Now, by their structure, planar diagrams give, in a confining theory, resonance behaviour[5] (with resonance widths proportional to N_f/N_c[6]). In this case, a (dual) interpretation of the parton model results in terms of resonances instead of partons should be possible. We can then ask if it is true that

$$\langle \sum_{Res} \sigma_{Res} \rangle \simeq \text{Parton Model.} \qquad (16)$$

When one looks into this problem for various QCD hard processes, one finds a little surprise: whereas for deep inelastic scattering and for e^+e^- annihilation one seems to be able to justify either Bloom–Gilman-type duality or Poggio–Quinn–Weinberg smoothing,[32] for lepton pair production a similar relation appears to be false. Without entering into details (see ref. 11), one finds that

$$\text{P. Model} \simeq |\langle A \rangle|^2 \overset{?}{\neq} \langle \sigma \rangle = \langle |A|^2 \rangle. \qquad (17)$$

In the presence of peaks, of course, $\langle |A|^2 \rangle > |\langle A \rangle|^2$ and one predicts that narrow peaks stick out of the DY predicted background (for instance, they do not have the famous $1/N_c = 1/3$ suppression from colour). This is not so unexpected, perhaps, but it is nice that the theory gives automatically this difference between the dual interpretation of various processes.

In the region of broad, strongly overlapping resonances, one will again expect the parton model prediction to be valid for the average physical cross-section.

The next question is how we do go from a parton to a resonance description of hadronic reactions. Unfortunately, we can only make, at present, some qualitative guesses about this difficult question.

§4. Low Energy, Spectroscopy

This is not the subject of my talk[33] and I will then touch it only very briefly in order to make a smooth transition into high energy, low p_T processes.

When we look at strong interaction phenomena at the large scale implied by energies of a GeV or so, we do not see any more the elementary constituents, but rather that complicated, coherent superposition of quarks, antiquarks and gluons which are the hadronic resonances and bound states. This is the regime where dual resonance models have been used in the past with reasonable success.

A suggestive connection between QCD and dual (string) theory has been proposed by several people and goes more or less as follows:

i) According to presently popular ideas,[34] confinement is a result of the exact local gauge invariance of the vacuum (as opposed, for instance, to a Higgs breaking situation).

ii) With a gauge invariant vacuum, only gauge invariant (colour singlet) states can propagate.

iii) The simplest gauge invariant states are formed by applying on the gauge invariant vacuum operators with a string-like structure, *e.g.*,

$$\bar{q}(x_1) T \exp\left(g \int_{P(x_1,\, x_2)} \mathrm{d}x_\mu A^\mu(x) \right) q(x_2) \sim \text{Open string, } x_1\, x_2$$

$$Tr T \exp\left(g \oint_P \mathrm{d}x_\mu A^\mu(x) \right) \sim \text{Closed string.}$$

The analogy goes a little further: if one uses the $1/N_c$ expansion of 't Hooft[5] the perturbative intermediate states of the leading diagrams have the same global colour structure as the states obtained by expanding in g those gauge invariant states. It can also be argued that, if confinement takes place, such states will represent (superpositions of) infinitely narrow hadrons.[5]

I would also like to mention a recent paper by Nambu[35] where further evidence for some possible connection of this sort has been given.

In Fig. 6 we show a table of correspondence between QCD-gauge invariant operators and strings corresponding to various hadrons. One interesting development in hadron spectroscopy, which can be studied this way, extending the original scheme of Rosner, is "baryonium" for which, however, I have to refer you to other parallel and plenary[33] sessions.

§5. Intermediate to High Energies (up to ISR?)

I shall be quite short on this part and refer you to the talk of Chan.[36] There has not been so much news in this area last year, whereas works of the previous year have mainly concentrated on baryon and baryonium, a subject discussed elsewhere.[33]

I shall limit myself to a brief account of the QCD topological expansion approach to the bare Reggeons and the Pomeron, also as an introduction to the next topic: Reggeon field theory.

As we move up in energy and we go to about the 10 GeV mark, particle production at low p_T starts to become important. In the QCD inspired string picture mentioned above,

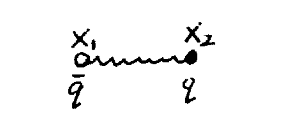
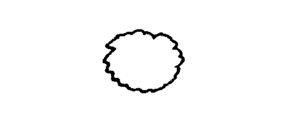
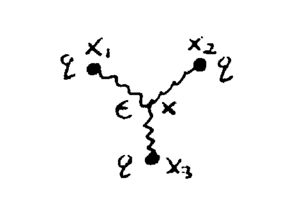
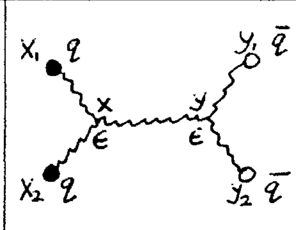
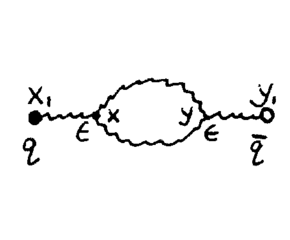
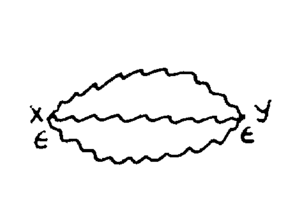
Hadron	Gauge invariant operator	String picture
M_2 = $q\bar{q}$ meson	$\bar{q}(x_1)\exp\left(g\int_{x_1}^{x_2} dx_\mu A^\mu\right) q(x_2)$	
M_0 = quarkless meson	$\mathrm{Tr}\exp\left(g\oint dx_\mu A^\mu\right)$	
B_3 = qqq baryon	$\varepsilon_{ijk}\left[\exp\left(g\int_x^{x_1} dx_\mu A^\mu\right)q(x_1)\right]_i \times \left[\exp\left(g\int_x^{x_2} dx_\mu A^\mu\right)q(x_2)\right]_j\left[\exp\left(g\int_x^{x_3} dx_\mu A^\mu\right)q(x_3)\right]_k$	
M_4^J = baryonium with $qq\bar{q}\bar{q}$ QN	$\varepsilon_{ijk}\varepsilon_{i'j'k'}\left[\bar{q}\exp\left(\int_y^{y_1}\right)\right]_i\left[\bar{q}\exp\left(\int_y^{y_2}\right)\right]_j \times \left[\exp\left(\int_x^{y}\right)\right]_k^{k'}\left[\exp\left(\int_{x_1}^{x}\right)q\right]^{j'}\left[\exp\left(\int_{x_2}^{x}\right)q\right]^{i'}$	
M_2^J = baryonium with $q\bar{q}$ QN	$\varepsilon_{ijk}\varepsilon_{i'j'k'}\left[\bar{q}\exp\left(\int_y^{y_1}\right)\right]_i \times \left[\exp\left(\int_x^{y}\right)\right]_j^{j'}\left[\exp\left(\int_x^{y}\right)\right]_k^{k'}\left[\exp\left(\int_{x_1}^{x}\right)q\right]^{i'}$	
M_0^J = quarkless baryonium	$\varepsilon_{ijk}\varepsilon_{i'j'k'}\times \left[\exp\left(\int_x^{y}\right)\right]_i^{i'}\left[\exp\left(\int_x^{y}\right)\right]_j^{j'}\left[\exp\left(\int_x^{y}\right)\right]_k^{k'}$	

Fig. 6. Simplest examples of correspondence between QCD gauge invariant operators and hadrons in the dual string picture.

we can say that long excited strings (heavy resonances) formed in the s channel break (decay). If the breaking up is a soft mechanism, as we believe, the fragments of the string will remember their location along the string itself, their ordering (see Fig. 7). As a result, a multiperipheral (MP) mechanism for hadronic production will follow, with approximate strong ordering, but also with cluster formation (clusters being $q\bar{q}$ resonances of a long enough lifetime).

Since the rate for string breaking (resonance decay) can be shown[16] to be proportional to N_f/N_c (at least for small widths), this parameter is seen to control the amount of clustering in the multiperipheral chain.

In any case, the shadow (overlap functions) of the production processes discussed above builds up (see Fig. 8a) an ideally mixed, OZI conserving, exchange degenerate Regge pole, the so-called planar bare Reggeon (*e.g.*, $\rho - A_2 - f - \omega$, $K^* - K^{**}$) which is the new collective excitation (quasi particle) replacing, in the sense of the old DHS duality, individual resonances.

Because of the simple fact that cutting a sheet

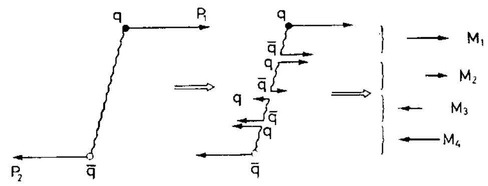


Fig. 7. Breaking of excited strings giving multiperipheral dynamics.

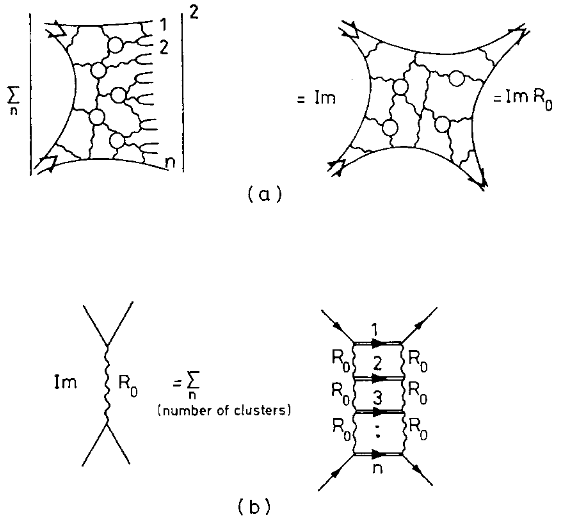


Fig. 8. (a) s channel content of the bare Reggeon (R_0). (b) Planar unitarity for R_0 within multiperipheral dynamics.

(plane) gives two sheets, the bare Reggeon satisfies a non-linear unitarity equation (see Fig. 8b) which is called planar unitarity. The unitarity sum is expected to be saturated in terms of the bare Reggeon itself with no AFS cut entering because of planarity.

By this kind of bootstrap, the strength of R_0 can be determined to some extent. The result is roughly:

$$\frac{g_R^2 N_f}{16\pi^2} \simeq .5(1+\bar{Y}_c/\bar{Y}_g)^{-1}, \qquad (18)$$

where $\bar{Y}_c(\bar{Y}_g)$ is the average cluster (gap) size in rapidity, both believed to be energy independent.

The last factor in eq. (18) is usually put to one by dual unitarization, topological expansion people, but it is actually important, both numerically and for over-all consistency with QCD.[16]

Numerically one gets

$$\sigma(\pi^+\pi^-)-\sigma(\pi^-\pi^-) \simeq \frac{100}{N_f}\alpha'(\alpha' s)^{-1/2}$$

$$\simeq \frac{40mb}{N_f}\left(\frac{s}{1Gev^2}\right) \qquad (19)$$

which is quite reasonable.

More theoretically, by the arguments made on the resonance lifetime, one finds that

$$\bar{Y}_c/\bar{Y}_g \sim N_c/N_f$$

and this, for $N_c \gg N_f$ gives

$$g_R^2 N_c/16\pi^2 \simeq 0(1)$$

which is an uncontroversial QCD result.[5,6]

Indeed, the two extreme cases $N_f/N_c \ll 1$ and $\gg 1$, correspond to two opposite pictures already considered in the past:

- $N_f/N_c \ll 1$ gives the narrow resonance picture of the ordinary dual loop expansion, but is very far from MP behaviour (no space for rapidity gaps).
- $N_f/N_c \gg 1$ corresponds to wide resonances and to direct multiperipheral production of pions, kaons, etc. It fits also with the Chew–Rosenzweig scheme of ***P***–f identify.[37]

Real life has $N_f/N_c \simeq O(1)$, and it is reassuring to find that, indeed, a model with MP production of clusters and with $\bar{Y}_c/\bar{Y}_g \simeq 1$ is more or less consistent with the data.[38]

Also the problem of f–***P*** identity can be clarified[6] by varying the parameter N_f/N_c and making contact in the relevant limits with dual perturbation theory,[39] which is known to give two separate vacuum trajectories.

The conclusion seems to be that it is quite easy to generate two vacuum singularities, but degeneracy of the non-leading one with the ρ, if experimentally estabished, will be a bit accidental in the topological expansion approach.

How does the transition to the bare Pomeron dominated regime take place? This is shown schematically in Fig. 9.

From resonance formation we go to a single multiperipheral chain with a wee valence quark being exchanged (this is also Feynman's picture explaining the $s^{-\beta}$ behaviour of quantum

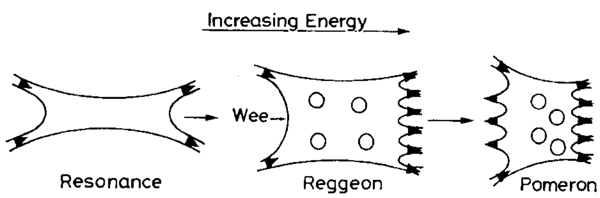


Fig. 9. Transition from the Resonance, to the Reggeon, to the Pomeron dominated regions.

number exchange cross-sections).

Increasing the energy further, it is increasingly difficult to keep the annihilating $q\bar{q}$ pair from emitting itself some mesons. When this takes place, vacuum exchange begins to be enhanced relative to quantum number exchange and this induces breakings of isospin degeneracy, ideal mixing, OZI rule, etc., in the region $t \lesssim 0$ (*vice versa* these properties can be shown[40] to become better and better as t grows large and positive).

The topological organizations of QCD (or dual) graphs tell us the precise relation between R_0 and the bare Pomeron $\boldsymbol{P}_0$ and allow one to define quite unambiguously the (bare) couplings of $\boldsymbol{P}_0$ to external particles, to Reggeons and to itself.

The basic picture for $\boldsymbol{P}_0$ is that of the shadow of a double MP chain (Fig. 9) with cluster production in each chain. The model is compatible with the ideas of Low and of Nussinov[41] on the Pomeron which also originated in QCD.

The basic properties of $\boldsymbol{P}_0$ coming from this picture (some of them are quite old results by now!) are the following:

If the two MP chains are independent:

i) $\langle n\rangle_{P_0} \xrightarrow[E\ \mathrm{large}]{} 2\langle n\rangle_{R_0}$; more generally, for the Feynman gas pressure:

$$p_{P_0}(z) \xrightarrow[E\ \mathrm{large}]{} 2p_{R_0}(z), \tag{20}$$

which has an amusing similarity to the relation found between quark and gluon jets for large N (eq. (13) with $C_F/C_A \to 1/2$).

ii) $$\alpha_{P_0}(0)=1 \tag{21}$$

iii) $$\alpha'_{P_0}(0)=1/2\,\alpha'_{R_0}(0). \tag{22}$$

These three relations will be somewhat modified by correlations among the two chains and, in particular, we expect $\alpha_{P_0}(0)>1$ and $\alpha'_{P_0}<\frac{1}{2}\alpha'_{R_0}$ to be correlated effects. In any case, the quantity $\Delta_0=\alpha_{P_0}(0)-1$ is of $O(1)$ in the $1/N$ expansion.

iv) One also finds that the triple Pomeron coupling at zero momentum transfer

$$g_{P_0P_0P_0}(0, 0, 0) \equiv g_P$$

is different from zero (strong coupling RFT) and that it is of $O(1/N)$.

All these quantities, R_0, $\boldsymbol{P}_0$, g_P belong to the $h=0$ topology of QCD graphs (sphere).

The amusing point is that higher topologies can be related[42] to RFT (Gribov) diagrams in a way consistent with s and t channel unitarity. This represents, to my knowledge, the only existing way of classifying all the diagrams of a normal field theory to obtain an RFT. The resulting RFT Lagrangian turns out to have the general structure[42]

$$L_{\mathrm{RFT}}^{\mathrm{TE}}=\bar{\phi}F(\phi/N, \bar{\phi}/N)\phi=\Delta_0\bar{\phi}\phi+\alpha'\bar{\phi}\boldsymbol{\partial}_t\phi + \frac{c}{N}\bar{\phi}(\phi+\bar{\phi})\phi+0(1/N^2) - 4\boldsymbol{P}_0\ \text{coupling}+\cdots. \tag{23}$$

An interesting question to ask is whether we should expect this RFT to be critical, *i. e.*, to have $\Delta_{\mathrm{ren}}=0$.

We know[43] that this happens for

$$\Delta_0=\Delta_c=g_P^2\times(\text{Known numbers}) \tag{24}$$

and we see therefore that, from our point of view, $\Delta_0=\Delta_c$ looks very accidental since $\Delta_0=O(1)$ and $\Delta_c=O(1/N^2)$.

For large N we certainly go in the supercritical direction $\Delta_0 \gg \Delta_c$ and in a simple but probably reliable model, Bishari has found[44]

$$\Delta_c \sim 0.004 \tag{25}$$

to be compared with $\Delta_0 \simeq 0.05$. We now turn to the discussion of recent results in RFT.

§6. Superasymptopia and RFT

The first question that comes to mind is: at what energies does RFT start to be relevant? The answer depends on the input parameters of RFT because they will decide when many $\boldsymbol{P}_0$ exchanges become important.

The general feeling is that, because of the small triple Pomeron coupling g_P and of the not-so-large Pomeron intercept, triple Pomeron iterations are not important yet in the ISR region. Since, however, the Pomeron coupling to the proton looks considerably bigger than g_P, other types of iterations (*e.g.*, eikonal) could already be relevant. Finally, if the coupling of $\boldsymbol{P}_0$ to a heavy nucleus goes like $A^{1/3}g_{PNN}$ (A=atomic number), then for scattering on nuclei, RFT could be already very relevant at present energies. This point about the coupling growing as $A^{1/3}$ looks, however, still controversial.[45]

Irrespectively now of relevance to present energies the conceptually important question comes of the expected behaviour of cross-sections, in particular of σ_{total}, at "∞" energies, In other words, how does a P_0 above one cure its own problem with the Froissart bound?

There has been considerable progress on these questions recently, part of which has been summarized by Le Bellac[46] in a parallel session. The results, however, although agreed upon by the majority of people working in this area, are still being challenged by White.[47] Let me go just a little into the basic points of super-critical ($\Delta_0>\Delta_c$) RFT.

a) Case of no transverse dimensions, $D_t=0$.

Even this case is controversial. It is actually believed[46] that settling this (apparently simple) case will also settle the dispute about the physical case $D_t=2$.

The S matrix for the forward elastic amplitude is written as

$$S=\langle 0|e^{-ifa}\, e^{-HY}\, e^{-iga^\dagger}|0\rangle=1+iT, \qquad (26)$$

where a and $a^\dagger$ are Pomeron destruction and creation operators, the two coherent states $\exp(-iga^\dagger)|0\rangle$, $\langle 0|\exp(-ifa)$ represent the (imaginary) couplings to the external legs of Fig. 10 and the $\exp(-HY)$ is the propagator blob (again of Fig. 10) in rapidity (imaginary time).

Inserting a complete set of eigenstates of H in Eq. (26) we see that the vacuum reproduces 1 and possible other eigenstates of H contribute to T, a term proportional to

$$\exp(+Y\Delta)=S^{+\Delta}$$

if $-\Delta$ is the eigenvalue.

Now H can be written down and the problem can be reduced[48] to the quantum mechanical problem of finding the levels of the corresponding Schrödinger equation. The potential looks quite different for $\Delta_0<0$ and $\Delta_0>0$ (see Figs. 11a and 11b).

The first case is non-controversial, but also uninteresting. For $\Delta_0>0$ the authors of ref. 48 find two low-lying levels, the vacuum and a second one which is above it by

$$-\Delta=\varepsilon\simeq\exp\left(-\frac{1}{2}\Delta_0^2/g_P^2\right) \qquad (27)$$

due to a tunnelling effect. A path integral formulation of such an effect, à la Polyakov, has also been obtained.[49]

As a result, one finds that, for $D_t=0$, there is no critical point at finite Δ_0. On the other hand, ε goes quickly to zero at increasing Δ_0.

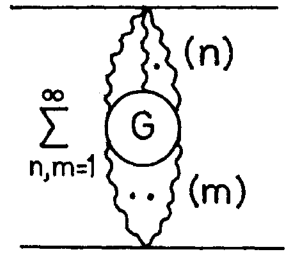


Fig. 10. RFT diagrams contributing to the elastic amplitude.

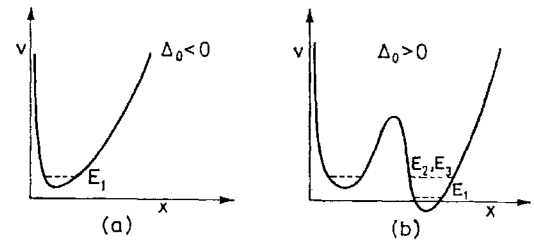


Fig. 11. Equivalent Schrödinger potential of ($D_t=0$) RFT. (a) $\alpha_P(0)<1$; (b) $\alpha_P(0)>1$.

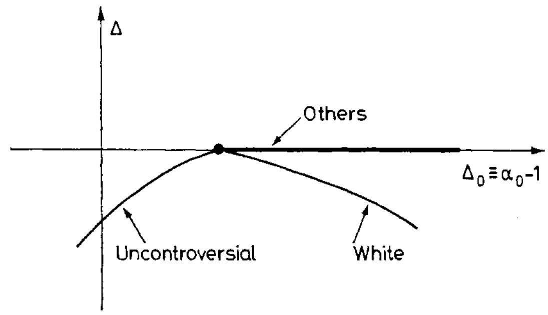


Fig. 12. Position of the output singularity (in $E=J-1$) as a function of Δ_0 . . . and of the theorist.

According to White,[47] instead, $\varepsilon\sim\Delta_0$ for Δ_0 large and positive and σ_T goes to zero faster and faster. I must confess, however, that I have not been able to understand his criticism.

$D_t=1,2$

The situation is summarized in Fig. 12, where the behaviour of the output **P** is given as a function of the input Δ_0. The region below and up to $\Delta_0=\Delta_c$ (under-critical and critical) is uncontroversial, but, as we argued, the real world probably lies above it. Here White appears to disagree again with the majority of the authors, whose model I shall now briefly describe.

The idea[50] is to neglect at first the small slope of the Pomeron. In this way, different points in impact parameter space (of dimensions D_t) are decoupled. Defining a lattice in such a space, we have to solve, at each lattice site, the $D_t=0$ problem discussed above.

For Δ_0 large, we can assume that only the two nearly degenerate states matter (the others are $0(\Delta_0)$ above). Hence, at each lattice site, we now have a system with two degrees of freedom, a "spin" variable which can be either up or down.

Coupling now the various lattice sites through α', we can ask whether the resulting spin model of RFT has a zero gap state. The answer is yes according to the authors of ref. 50. They find a zero gap state $|1\rangle(E_1=0)$ with

$$\langle 1|a^+(y, b)|0\rangle \simeq \frac{\Delta_0}{g_p}\theta(vy-b), \quad (28)$$

where b is the impact parameter, y the rapidity and v a calculable parameter. Equation (28) is equivalent to a grey expanding disc giving $\sigma_T \sim \log^2 s$, as in the Froissart bound.

For $D_t=1$ even a soliton, path integral formulation of the phenomenon can be given.[51]

Various authors[52] have been able to compute multiparticle distributions in the above spin model and have found a sort of modified geometrical scaling behaviour for the n particle density $\rho^{(n)}$:

$$\rho^{(n)}(y_i) \sim (Y^2)^n F_n(y_i/Y) \quad (29)$$

to be compared with the critical point behaviour:

$$\rho^{(n)}(y_i) \sim Y^{n\eta}\tilde{F}_n(y_i/Y). \quad (30)$$

Both eqs. (29) and (30) are compatible with KNO scaling.

Finally, exclusive and inclusive diffractive production are found[53] to be damped enough for s channel unitarity constraints to be fulfilled.

In conclusion, the spin model looks like a very consistent scheme for supercritical RFT and is able to predict rising cross-sections. If it will pass, as I feel it should, the further test of t channel unitarity, the remaining theoretical objections to it should probably be dismissed.

§7. Soft *vs* Hard Hadron Physics

7.1. *Phenomenological analogies*

Of course the jet structure is common to both types of physics (large angle jets *vs* forward and backward jets). This has led to the speculation of complete jet universality between soft and hard processes.

Data suggest an approximate validity of these predictions, *e. g.*,

a) Average multiplicities are expected to be related according to[54]

$$\bar{n}_{e^+e^-}(E):\bar{n}_{\pi P}(\alpha E):\bar{n}_{\bar{p}p_{\text{ann}}}(\beta E)=1:2:3, \quad (31)$$

where α, β are more or less predictable unmbers[55] accounting for the fact that in e^+e^- all the energy goes to a single $q\bar{q}$ pair whereas in other processes it gets distributed among various $q\bar{q}$ pairs (*e.g.*, three pairs in $\bar{p}p$ annihilation).

Although the data now clearly disagree[56] with absolutely universal behoviour of the type predicted in ref. 26, they seem to be roughly in agreement with the prediction (31).

There are other predictions on correlations, in particular on Bose–Einstein interference,[57] which also look as if they have experimental support.[58]

b) *There have been claims*[59] that fragmentation functions in e^+e^- and low p_T hadron physics scale and are roughly consistent with each other, although, in a contribution to this Conference,[60] the opposite statement has been made.

Notice also that the average $\langle q_T\rangle$ relative to the jet axis has been reported[61] to be the same as that of low p_T physics (~ 300 MeV).

7.2. *Theoretical QCD expectations*

It looks that, in QCD, the two types of jets should differ drastically at a closer analysis. The analogy should be only superficial and/or just restricted to low values of Q^2. I shall try to argue now that this is indeed our theoretical expectation.

a) We can quote first results from two-dimensional QCD,[62] where one finds that, whereas fragmentation functions are universal within hard processes (as we argued to be also the case in four dimensions), they differ in soft hadron-hadron collisions.

In the actual four-dimensional case there are further reasons to doubt universality.

b) Firstly, low p_T fragmentation functions should scale, to first approximation, if QCD gives, after confinement, Regge poles interpolating its bound states. Of course there will be scaling violations associated with absorptive corrections (Regge cuts) but these are non-planar effects down at least $O(1/N^2)$.

On the contrary we have seen that, in hard

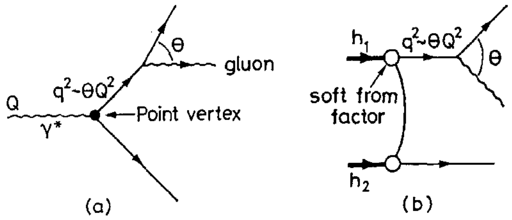


Fig. 13. Producing a large p_T in (a) hard processes (*e. g.*, $e^+e^- \to$hadrons); (b) soft processes (*e. g.*, $K\bar{K} \to$pions).

processes, scaling violations in fragmentation functions are there already (if not only) in the simplest planar diagrams. At some Q^2 they will start being large and $O(1)$ (in terms of a $1/N$ expansion).

c) A very much related difference is the fact that low p_T physics is strongly damped in p_T: roughly speaking $\langle p_T \rangle = 300$ MeV and constant.

In hard jet physics we expect the average p_T relative to the jet axis to increase essentially linearly with energy (same for the p_T of lepton pairs). In other words, hard jets have finite angular spread and *not* finite p_T.

As I anticipated, b) and c) have a common origin in the rather hard nature of QCD (or of any other field theory with dimensionless coupling constant). The mechanism giving a different behaviour in the two regimes is sketched in Figs. 13a and 13b. In Fig. 13a it is "easy" to send the intermediate quark line much off shell (which is necessary in order to produce a large p_T relative to the jet axis) because of the point-like coupling of the photon (a similar effect occurs also in large p_T hadron hadron collisions). On the contrary, in Fig. 13b, the same process is killed (exponentially?) by the wave function of our composite system which does not like to have partons far from the mass shell.

In other words, once a hard process is generated, it is only a little extra price ($\alpha_s(Q^2)/\pi \sim 0.1$) to produce an even harder one (*e. g.*, large q_T relative to the axis of a jet of large p_T); to do it instead the first time (*i. e.*, large p_T relative to the beam direction) is much harder.*

How do we then explain the apparent analogies between soft and hard hadron physics? Well, up to present values of s and Q^2, very large scaling violations are neither predicted nor seen: universality could very well be an approximate and temporary property before we move to harder and harder processes.

In any case, this looks to be a distinct prediction of QCD, as opposed to softer theories like Preparata's bag model or conventional dual theories, and should be checked at the energy of next generation machines.

I would like to conclude with a speculation on what could be instead a real common denominator to soft and hard hadronic phenomena. The guess is that such will be the concept of "order" in hadron physics be it soft or hard.

It looks now very plausible that, in the axial gauge of QCD, a small set of diagrams, having the simplest topology, describe to great accuracy both hard and soft processes, both lepton and hadron induced reactions. The manifestations of this fact would be numerous, *e. g.*,

a) The orderly, coherent structure of hadronic production, to be detected perhaps by sensitive interference effects of the Bose–Einstein type.

b) The simple correlation of momenta and internal quantum numbers in jets.

c) The dominance of resonance production (clusters) in low-energy channels.

d) The OZI and EXD regularities together with the pattern of their breaking.

Actually, for low p_T hadron physics, Chew and co-workers have been recently setting up an ambitious S matrix program having the concept of order as its basic starting point. This, I feel, could extend to hard processes as well and provide* a distinctive feature of strong processes as opposed to other types of interactions.

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* An interesting piece of data which seems to confirm such QCD prediction in the observation[61] of a larger $\langle q_T \rangle$ in the forward-backward fragments, where a large p_T trigger is used. The QCD expectation is that $\langle q_T^2 \rangle \sim O(\alpha_s(p_T^2) p_T^2)$.

* It could also provide a useful hint to confiners. For instance, instanton effects are expected to turn off in the planar limit (they are $O(e^{-N})$) and, therefore, if the planar theory already confines, as we argued to be the case, instantons would have little to do with that phenomenon.

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NOTE ADDED IN PROOF

Soon after the Conference, the following papers have appeared extending the analysis of QCD jets to the "semi-hard" region discussed in Section 3.2:

G. Curci and M. Greco: CERN Preprint TH–2526 (1978); R. K. Ellis and R. Petronzio: CERN Preprint TH–2571 (1978); K. Konishi, A. Ukowa and G. Veneziano: CERN Preprint TH–2577 (1978).

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P 3b Dynamics of High Energy Reactions

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Applications of the theory of quantum chromodynamics (QCD) with asymptotic freedom to processes involving large momentum transfers are examined. The theory describes correctly many features of the lepton initiated processes eN, μN, νN, $\bar{\nu}N$ as well as large-mass muon-pair production *and* the production of mesons and jets at large $p_\perp$ in hadron-hadron collisions. The preliminary conclusion is that QCD might well be the correct theory behind all these phenomena, although a definitive test has not yet been made.

§I. Introduction

During the last several years, a new framework to describe strong interaction physics has emerged: quantum chromodynamics (QCD). It is the simplest field theory which incorporates a color-dependent force among the quarks. These forces are generated by the exchange of colored vector gluons which are coupled to the quarks (and to each other) in a gauge-invariant manner. The theory is closely related to the most successful quantum field theory: QED. The only (but very important) difference is the gauge group involved. QED is an Abelian gauge theory (the photons do not couple to each other); QCD is a non-Abelian gauge theory—gauge group SU_3 (color). The gluons carry color and thus couple to each other.

Although the theory is well defined, precisely what it predicts is not yet clearly known. For example, it is not known if the theory actually confines quarks and gluons within hadrons nor has the spectrum of hadron states been calculated. At present, the mathematical complexities are still too great. However, at very high energy or momentum transfer Q, the theory is asymptotically free; the effective coupling between quarks and gluons decreases toward zero with increasing Q^2. As emphasized by Politzer,[1,2] this permits calculation of those parts of a process involving high Q^2 by the use of perturbation theory. Yet most real processes involve both high *and* low Q^2 together and precisely how to separate these parts is just becoming understood.

In this talk, I will examine many of the present day applications of QCD to processes involving large momentum transfers. Some of these applications are rather crude and involve ideas that are somewhat phenomenological in nature. Nevertheless, comparisons with data are quite encouraging, although many of the most dramatic (and definitive) tests are yet to come. The theory describes correctly many features of the lepton initiated processes eN, μN, νN, $\bar{\nu}N$ as well as large-mass muon-pair production *and* the production of mesons and jets at large $p_\perp$ in hadron-hadron collisions. The preliminary conclusion is that quantum chromodynamics might well be the correct theory behind all these phenomena.

§II. The Effective Coupling $\alpha_s(Q^2)$

The theory of QCD does not produce inclusive cross sections that "scale." One cannot use dimensional counting arguments to determine the behavior of cross sections (at intermediate values of Q^2). This is because the theory has an intrinsic "scale" or mass parameter Λ that is generated as a result of the *interaction* between quarks and gluons. These interactions result in an effective strong interaction coupling, $\alpha_s(Q^2)$, that decreases logarithmically with increasing Q^2, where Q is some characteristic momentum in a collision.

In the theory of QED, it is well known that the physical coupling e, defined by the large distance (small Q^2) behavior, is smaller than the effective coupling e_{eff} one would measure at small distances (large Q^2). This is due to vacuum polarization effects that shield the bare charge. In lowest order perturbation theory, the vacuum polarization contribution shown in Fig. 1a gives