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COLLISIONS OF CONSTITUENT QUARKS AT COLLIDER ENERGIES *

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ABSTRACT

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1 Constituent Quarks versus Quark-Partons

How many quarks are there in a proton? Even a fair number of little children will provide the quick answer of three, not to mention a large number of adult physicists. I asked this of Dick Taylor during this meeting, just to see how he would answer. After all if anyone should know it would be him. Dick is a savvy fellow and immediately smelled a trap, and gave a quite correct albeit wimpy reply... "Well, it depends, doesn't it?" Yes it does, because if the proton is at high momentum and deep inelastic scattering is used to count the quarks, one gets

$$\int_0^1 \frac{dx}{x} F_2(x) = \sum_i e_i^2 = \langle e^2 \rangle N \cong \frac{2}{9} N .$$

$$\tag{1}$$

If one integrates over all x, corresponding to an infinite-momentum proton, the number is infinite. And even for a practical cutoff of $x > 10^{-4}$, the integral gives a value of N of about 30, and growing rapidly with decreasing x.

So why the three? The three is the number of constituent quarks, while the infinity or whatever is the number of quark-partons, And what are the differences?

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- 1. Constituent quarks are extended objects, while the quark-partons are pointlike.
- 2. The mass assigned to constituent quarks is about 350 MeV for u and d constituent quarks, and about 500 MeV for s. On the other hand the quark-partons, or "current-quarks", are assigned masses of about 4, 7, and 150 MeV respectively.
- 3. The constituent quarks (and/or antiquarks) are "lightly bound" to form the hadrons; quark-partons are relevant at short distances and are "asymptotically free".
- 4. The effective coupling to gluons of constituent quarks is not large, despite the relatively low momentum scale, or large distance scale. The couplings of quark-partons to each other and to gluon-partons is perturbative, with the strength controlled by the (small) running coupling constant of QCD.
- 5. Probably the constituent quarks are "small", with a radius of order 0.2-0.3 fermi. This accounts for the success of the additive quark model in interpreting the nature of the high-energy total cross sections for hadron-hadron collisions.

The evidence for the constituent-quark predates that for the quark-parton; it mostly resides in the successful interpretation of hadron spectroscopy via SU(6) and the ideas of the quark model. But there are many unanswered questions regarding the internal structure of these constituent quarks and consequently of the hadrons themselves.

In these lectures we concentrate on high energy collisions of hadrons with each other, or with an electron, and consequently on the "light-cone" structure of the hadrons. That is, the hadrons we deal with will be extreme relativistic: pancakes of partonic matter whose constituents' internal motions are frozen in place during the course of a collision.

An enormous amount of effort has gone into determining the longitudinal momentum distributions of these parton constituents. But there still is no accepted simple picture of the transverse impact-plane structure of the hadron. The default viewpoint is that it is just an uncorrelated gas (Fig. 1a). But perhaps the constituent-quark model alluded to above would suggest a "nuclear" model, with most of the parton distribution concentrated inside the constituent quarks. Or maybe the hadron is stringy with most of the partons in the string. It intuitively seems that phenomenology should be sensitive to these options. But it is not so easy, and in fact not too much has been accomplished, either experimental or theoretical, so far.

Part of the problem is that the transverse structure of ordinary mesons and especially of baryons is sure to be complicated, just because more than one constituent quark is present. However with the advent of heavy flavors, there are in principle simpler systems containing at most one constituent quark. Onium is perhaps the



Figure 1: What does the proton look like? (a) Uncorrelated parton gas; (b) constituent quarks; (c) strings.

cleanest; perturbative QCD should nearly suffice to determine its structure-although even here there are complications we will note later on. The next simplest systems are heavy mesons like the B, which contain only one constituent quark. A high energy B in heavy-quark effective theory is just a high-energy constituent quark plus minor perturbatively calculable effects due to the spectator \overline{b} quark going along for the ride.

So it would be nice to have beams of B's and Υ 's instead of pions and nucleons; the dynamics would be much simpler. While that is not so practical to realize, it turns out that there are good substitutes: the virtual photons of e-p collisions. We will return to this later.

2 Pion Clouds and Hadron Structure

Constituent quarks should have pion clouds surrounding them. This follows from the inference that their sizable mass of 350 MeV is due to chiral symmetry breaking. Then the coupling of pion to quark follows from the Goldberger-Treiman relation [1], now applied to the constituent quark.

$$2m_Q \left(\frac{g_A}{g_V}\right) = \sqrt{2} f_\pi g_{\pi QQ} \ . \tag{2}$$

The mathematics of this description is conveniently given by the σ -model. It is a precise strong-interaction analogue of the Higgs Lagrangian. Four phenomenological fields are introduced: three are the Nambu-Goldstone pion modes (analogous to longitudinal W's and Z's) and the fourth is the σ , with vacuum quantum numbers and a nonvanishing vacuum-expectation-value (analogous to the Higgs field). The σ -model is used as an effective field theory, valid only at relatively low energies. Manohar and Georgi argue [2] that "low energy" means "under 1 GeV", with the scale given by $(4\pi)f_{\pi}$, with f_{π} the pion decay constant, already encountered above in the Goldberger-Treiman relation. However there are some differences. Here the pion coupling is quite strong:

$$\left(\frac{g^2}{4\pi}\right)_{\pi QQ} \approx 2 \tag{3}$$

while the corresponding coupling of Higgs system even to top quark is only about 1/30. The quartic coupling, λ , between σ 's and pions is likewise quite large

$$\lambda \approx 20 \tag{4}$$

as estimated from the "decay width" $\Gamma_{\sigma}/m_{\sigma} = 1/2$. The Higgs boson would have to have a mass in the TeV region to attain such a large coupling.

So we may anticipate that pion clouds of diameter 1–2 fermi will surround constituent quarks. How does this affect the distribution of partons associated with a constituent quark? Do most of the (primordial, $Q^2 \leq 3 - 10 \text{ GeV}^2$) partons reside in the cloud?

At present my preference is to answer this question yes. In the last few years I tended to assume no; the partons were inside the constituent quark (cf. Figs. 1b and 1c). However this puts a lot of transverse momentum into the parton wave functions, even a large mean transverse momentum. The data on primordial transverse momentum comes from deep- inelastic scattering, especially the limits on $R = \sigma_L/\sigma_T$, and from analyses of the Drell-Yan processes of hadroproduction of dileptons. Quite a bit is allowed, but I could not scratch out consistent numbers nevertheless [3].

It seems more satisfactory to keep the partons on the outside of the constituent quark, and to define the region where the constituent quark *is* as the region where the chiral vacuum condensate *is not*. As an idealization, only the valence parton and its perturbative (Altarelli-Parisi) evolutes would be inside. On the outside are all the ocean partons, immersed in the ocean of chiral condensate (Fig. 2). This has the advantage that the mean momentum of these outsider-partons need not be controlled by the size of the constituent quark. Because the wave functions of these ocean partons are excluded from the core of the constituent quark, they will have high-momentum tails, however.

The gluon field of the valence quark-parton is strongest in the core; it does not peacefully coexist with the chiral condensate. So the classic picture of stringy fluxtubes between the valence degrees of freedom fits easily into this viewpoint as well.

What about onium? Is there a pion cloud around it? The answer is yes, there is some. Upsilons, for example, can mix with $B\overline{B}$ continuum states without the cost of a big energy denominator. And the $B\overline{B}$ systems can exchange pion pairs. So at large distances there will be a pion cloud. And at sufficiently high momenta this pion cloud will have enough momentum density (GeV/square fermi) that long-range peripheral collisions will occur.



Figure 2: Are the primordial partons outside the constituent quarks? (a) Constituent quarks plus cloud; (b) strings plus cloud.

3 Implications for Collision Phenomena

Ideally it would be nice to have beams of Υ 's to collide against each other at high energies. In fact Mueller uses nothing else when doing his theoretical studies of hard diffraction. And the next best things would be beams of B mesons, containing only one constituent quark per meson.

It will be a long time before such beams exist in real life. But it turns out that there is a good substitute: virtual photons. These exist in abundance at HERA.

Why is this so? Consider a virtual-photon proton collision at very high cms energy. An interaction will occur when, upstream of the collision point, the virtual photon fluctuates into a quark-antiquark pair. At birth this pair consists of two pointlike partons; typically their transverse separation when reaching the collision point will be dictated by the uncertainty principle and be of order $(Q^2)^{-1/2}$. This is the case when the partition of longitudinal momentum to quark and antiquark is typical, of order unity. Under these circumstances the virtual-photon system at arrival is a small color dipole which interacts perturbatively with the nucleon. This is not so different from idealized onium, which again can—for sufficiently heavy quarks—be regarded in a similar way as a small color dipole.

However most of the deep-inelastic interactions which are observed occur when the momentum partition of the quark and antiquark is not typical, but when one carries almost all the momentum and the other a momentum fraction of order $(1 \text{ GeV}^2/Q^2)$. When this configuration occurs, the alignment of the $q - \bar{q}$ system in its rest frame, relative to the collision axis, is sufficient to render the transverse momenta of q and \bar{q} less than 1 GeV. The uncertainty principle then allows them to be separated (transversely) at arrival to the target by a distance of order the hadron size. Consequently there is enough time for the slow quark or antiquark to "dress" itself with strings, pion clouds, and/or other nonperturbative stuff. In other words the virtual photon system becomes hadronlike, and interacts strongly with the nucleon in the collision.

In fact the system may be similar to the B meson; the fast quark is analogous to the b, perturbative and relatively passive, while the slow quark (or antiquark) is analogous to the constituent quark in the B: structured and strongly interacting [4].

The distinction between the two cases can be defined event-by-event by the finalstate properties. In the former case there will in the photon-fragmentation region of the lego plot be a pair of jets each having p_t of order $(Q^2)^{-1/2}$, with the pair carrying almost all the momentum of the incident virtual photon. In the latter "aligned-jet" case, there will be essentially a featureless, jetless lego plot.

Experimentally the "aligned-jet" case is what is typically found at HERA and fixed-target experiments, as well as what is anticipated just from naive parton-model ideology (and Monte-Carlos). In addition the latter mechanism predicts that at small x there will be nuclear shadowing ($A^{2/3}$ behavior) because the virtual-photon configuration is hadronlike.

So we may consider the typical small-x deep-inelastic process as very analogous to *B*-nucleon interactions at high energy: one constituent quark is absorbed onto the target nucleon. And with this preliminary out of the way, we can get down to the question of whether there is any evidence that pion clouds around such constituent quarks are of importance in the dynamics. One possible piece of evidence is the discrepancy in the Gottfried sum rule between observations and naive expectations. The sum rule is

$$\int_{0}^{1} \frac{dx}{x} \left[F_{p}(x) - F_{n}(x) \right] = \frac{1}{3} \int_{0}^{1} dx \left[u(x) + \overline{u}(x) - d(x) - \overline{d}(x) \right]$$
(5)
$$= \frac{1}{3} + \frac{2}{3} \int_{0}^{1} dx \left[\overline{u}(x) - \overline{d}(x) \right] .$$

The difference between the experimental value of 0.22 and the simple expectation of 0.33 is evidence that there are more \overline{d} quarks surrounding the proton than \overline{u} . Since there are more u quarks than d in the proton, this can be interpreted as more π^+ than π^- cloud in the proton, because the extra \overline{d} 's are made in the virtual transition $u \to d + \pi^+ = ud\overline{d}$. Likewise the analysis of the spin sum rules [5] leads to a value of Δs less than zero; *i.e.* the strange-quark cloud is polarized oppositely to the u quarks in the proton which carry the spin. This can be interpreted [6] in terms of the spin-flip occurring when the kaon cloud is made via $u \to s + K^+$.

However it would be much nicer if these features, in particular the density of partons in the impact plane, could be seen directly. Usually this is not possible, because in the typical hard-collision experiments the transverse distributions are "integrated out". But an exception to the rule is in double hard processes. These are the class of events where two hard collisions occur in the same event, leading for example to two pairs of isolated coplanar dijets. If the longitudinal fractions of the initial-state partons which initiate the collisions are small enough, then the joint cross section is just

$$\frac{d\sigma(1,2)}{d\Gamma_1 d\Gamma_2} = \frac{d\sigma(1)}{d\Gamma_1} \frac{d\sigma(2)}{d\Gamma_2} \frac{1}{\sigma_{\text{eff}}} .$$
(6)

Here $d\Gamma_1$ and $d\Gamma_2$ are the final phase-space elements for the produced dijet pairs 1 and 2 respectively. The factor σ_{eff} is just the measure of correlation of the primary partons in the impact plane; it has the dimensions of a cross section. If σ_{eff} is small there will be a lot of impact-plane correlation.

Experimentally, the situation is confused. At the CERN ISR, the AFS collaboration first searched for double parton interactions and claimed [7] a small value for σ_{eff} . Then, at UA2, a lower bound was set [8], while at the TeVatron the CDF group marginally claims [9] a small value of σ_{eff} . Theoretically, there has been occasional work on the subject [10], but in my opinion there has yet not appeared comprehensive theoretical work either.

Stimulated by some work in progress by the D0 collaboration [11] at Fermilab, I made a short study of the subject, concentrating on impact-plane correlations. (In practice, this is not enough; careful attention must be paid to correlations in the longitudinal fractions—especially those dictated by energy conservation.) I expected to find a strong sensitivity to clustering of partons around constituent quarks, but did not find as much as expected. Some of the reason can be seen just from considering the simple example of B-B scattering where in the constituent quark model one has (cf. Fig. 3)



Figure 3: Double-parton scattering geometry for a $B-\overline{B}$ collision.

$$\frac{1}{\sigma_{\text{eff}}} = \int d^2 b_1 \, d^2 b_2 \, d^2 B \, f_A(b_1, b_2) \, f_B(B - b_1, B - b_2) \,. \tag{7}$$

Here $f_A(b_1, b_2)$ is the joint distribution function for finding partons 1 and 2 at impact parameters b_1 and b_2 in projectile A. Clearly σ_{eff} senses the size of the constituent quark. But so does the total *B-B* cross section, which by the additive quark model would be expected to be about 4 mb. So how much insight can be attained by study of the double-parton processes is not yet clear. But it is probably worthwhile to investigate more, both theoretically and experimentally.

Now let us turn in more detail to the dynamics of generic, soft high energy collisions. At the simplest descriptive level, these are expressed, for better or worse, in terms of meson exchanges. For B-B scattering at small impact parameters, we expect one-gluon exchange to be predominant, especially between the *b*-quarks. At the largest impact parameters it must be a pion which is exchanged, because it has the smallest mass. The typical amplitudes for single gluon exchange have the form

$$A \sim \alpha_s \left(\frac{s}{t}\right) \tag{8}$$

or

$$\frac{d\sigma}{dt} \sim \frac{\alpha_s^2}{t^2} \tag{9}$$

while for single pion exchange between constituent quarks we have

$$A \sim \left(\frac{g^2}{4\pi}\right) \ \frac{\langle \sigma \cdot q \rangle \langle \sigma \cdot q \rangle}{(t - m_\pi^2)} \sim \left(\frac{g^2}{4\pi}\right) \ \frac{t}{(t - m_\pi^2)} \tag{10}$$

$$\frac{d\sigma}{dt} \sim \frac{1}{s^2} \left(\frac{g^2}{4\pi}\right)^2 F(t) . \tag{11}$$

These exchanges, however, evolve with increasing energy into "multi- peripheral" or ladder exchanges. This occurs because, as in QED, exchange in one frame (Coulomb photon) becomes emission of a quantum (Weiszacker-Williams virtual photon) in another, and the emitted object can branch into pairs of other objects. It is this potential for branching which creates the ladder. (This is important in QED as well; muon interactions at very high energies feel the electron-positron "conversions" in the photon-exchange ladder.)

These "ladder" exchanges extend the range of the force and modify the energy dependence of the interaction, in particular increasing the strength of the lowest order process at high energies. This change in energy dependence is dubbed "Reggeization". If the net energy dependence of the forward scattering amplitude turns out to be a pure power, this is called Regge-pole exchange, for reasons to be discussed later. If the dependence is not a pure power, one says there is a Regge-cut.

Note that the ladder exchange refers to the forward scattering amplitude, not the inelastic amplitudes that build the imaginary part of the forward amplitude (total cross section). So the primitive ladder amplitude is comprised typically (but not always) of exchange of the particle and its antiparticle. The cases relevant to high energy hadron scattering are perturbative two-gluon exchange (the hard, perturbative, QCD Pomeron), "soft" Pomeron exchange, and Reggeized non-singlet meson exchange, e.g. ρ exchange. The QCD "hard" Pomeron is reasonably calculable from perturbation theory. Rho-exchange becomes a pure Regge-pole process; it is well

or

documented experimentally in pion charge exchange scattering and neutral-kaon regeneration, and even in non-singlet deep-inelastic scattering structure functions. The candidate exchange for the ρ Reggeon is evidently the (constituent) quark-antiquark pair from which the meson is built.

The soft-Pomeron exchange is, by definition, what contributes to the hadronhadron total cross sections at high energy. The quanta that build its "ladder" (if there is one) are the least certain. The most popular choice is a pair of non-perturbative gluons [12]. However if collisions of pion clouds surrounding constituent quarks are an important mechanism at high energy, then it is at least arguable that exchange of pion pairs, σ 's, and constituent-quark pairs may work together to build the soft Pomeron, with gluons relatively unimportant. This is the point of view taken here. But it is a speculative one.



Figure 4: Ladders and their origin: (a) Multigluon production amplitude, whose square (b) builds the ladder for the QCD "hard Pomeron"; (c) Meson-exchange ladder which builds the ρ Regge trajectory; (d) A possible ladder for building the "soft Pomeron".

To summarize, the structure of these ladders (cf. Fig. 4) are supposed to have the basic properties exhibited in the following table:

In the next section we shall explore in somewhat more detail what these entries mean, and the relation between the ladder structures and Reggeon theory.

4 Basics of Regge Behavior

The theory of Regge poles is a quite dormant topic. It does not seem to be taught very much any more. In addition there is often found an attitude that the subject is

	ρ Reggeon	Hard QCD Pomeron	Soft Pomeron
Composition	$(\text{constituent})\overline{Q} - \overline{\overline{Q}}$	Gluon pair	Gluons?
(ladder sides)			Pion pairs?
J-plane structure	Pole	Cut	Pole???
s-dependence	$s^{-0.45}$ (for $\Delta\sigma$)	$s^{0.4}(\log s)^{-3/2}$	$s^{1.08}$
Coupling	Constituent quarks	Pointlike	Constituent quarks
			color dipoles
Observed?	Cross section differences	HERA F_2 ????	$\sigma_{ m total}$
	$d\sigma/dt$		
	$F_p^2 - F_n^2$		
Regge-trajectory	1 GeV^{-2}	Predicted to be	$0.25 { m ~GeV^{-2}}$
slope		very small	·

obsolete, because it is identified so strongly with the prequark, pre-parton era of the S-matrix, dispersion-relations approach to strong interactions.

This point of view is just plain wrong. The Chew, Frautschi, Regge, et al. description of high energy behavior in terms of singularities in the complex angular momentum plane is completely general (This is true of most of the dispersion relations, too.) In addition, their conjecture that high energy behavior of two-body scattering amplitudes might be describable in terms of moving poles in the J-plane is beautifully verified in the case of non-singlet meson exchanges. And the basic technique of Watson-Sommerfeld transform should be a standard part of the training in theoretical particle physics.

Part of the problem may be that the literature on the subject is not too easily accessible. For these lectures I searched briefly, but largely in vain, for a succinct, easily accessible source. There are books by Collins and Squires [13] and by Newton [14] that help, but they are not too easy. An early paper by Frautschi, Gell-Mann, and Zachariasen [15] reads well. I am sure there are other good sources, and that it was just lack of time/effort on my part which is to blame. I will appreciate learning of the favorite source material of others.

The most direct approach to the subject is a la Regge, using an analysis of Bethe-Salpeter ladder equations. This will not be attempted here; one must learn a technology of P's and Q's (properties of Legendre functions P_l and Q_l), and there is neither the space nor lecturer expertise to do it full justice here. Instead, two more schematic approaches will be used, the multiperipheral and the Feynman-Van Hove.



Figure 5: Building a soft Pomeron: (a) Lowest order pion exchange; (b)Emission of one σ ; (c) Multiple σ emission.

The multiperipheral approach, developed by Fubini and collaborators [16], builds the Reggeon by simply summing the individual contributions to the ladder for the imaginary part of the amplitude (inelastic cross section). We exhibit it for our simple model of the soft Pomeron: pion exchange. The lowest order amplitude is shown in Fig. 5a. The cross section, when cut off at large t, falls with squared cms energy slike s^{-2} :

$$\sigma_0 \sim g^4 s^{-2} . \tag{12}$$

We will take the rungs of the ladder to be the J = 0, spinless, σ -particle of the σ model (the Higgs particle of the strong interactions!), which phenomenologically is a very broad s-wave $I = 0 \pi - \pi$ resonance at about 700 MeV. The cross section for production of one σ is, for large enough s,

$$d\sigma_1 \sim (s_1 s_2)^{-2} g^4 G^2 d\Gamma_1 \tag{13}$$

where s_1 and s_2 are the squared subenergies for the process $q + \pi \rightarrow q + \sigma$ (Fig. 5b), and where $d\Gamma_1$ is the phase space element for the σ , d^3p/E . The cross section for production of $n \sigma$'s follows in the same way:

$$d\sigma_1 \sim (s_1 s_2 \dots s_{n+1})^{-2} g^4 G^2 d\Gamma_1 \dots d\Gamma_n .$$
⁽¹⁴⁾

A simplification occurs when the coupling constant G is small, so that the density of produced σ 's in the lego plot is low. The sum of logs of the subenergies just is the sum of rapidity separations of the σ 's in the lego plot, which adds up to the total energy (Fig. 6). Now one has to integrate over the phase space $d^2p_td\eta$ of the produced σ 's. We only integrate over low p_t —either the cutoff is automatic or else



Figure 6: Lego Plots for the processes of Fig. 5.

it is unrealistic to extend the production model to large p_t . The more important integration is over the (pseudo)rapidities, which are ordered:

$$\int_0^{\ell_n s} d\eta_1 \int_{\eta_1}^{\ell_n s} d\eta_1 \dots \int_{\eta_{n-1}}^{\ell_n s} d\eta_n = \frac{(\ell_n s)^n}{n!} .$$
 (15)

Then summing things up gives

$$\sigma = g^4 s^{-2} \sum_{n=0}^{\infty} \frac{1}{n!} \left[\int \frac{d^2 p_t}{(2\pi)^3} G^2(p_t) \right]^n .$$
 (16)

There are several important and quite generic features to note here. First, each of the individual contributions had, up to logarithms of s, the same s^{-2} power behavior. But after the summation the exponent of s, which we shall call α , changed, and that the change was positive:

$$\tau \sim s^{(\alpha-1)} \tag{17}$$

with

$$(\alpha - 1) = -2 + \int \frac{d^2 p_t}{(2\pi)^2} G^2(p_t) .$$
(18)

Note also that the mean number of ladder "rungs", *i.e.* the mean number of produced σ 's, is

$$\langle n \rangle = (\ell n s) \int \frac{d^2 p_t}{(2\pi)^3} G^2(p_t)$$
(19)

and that the distribution around the mean is Poisson. This also gives the mean number of ladder "rungs", or σ 's, per unit rapidity as

$$\frac{dN_{\sigma}}{d\eta} = \int \frac{d^2 p_t}{(2\pi)^3} G^2(p_t) = \alpha + 1 .$$
 (20)

While the model sketched here is rough and simplistic, there is a generality associated with these conclusions. For the soft Pomeron, we would infer that, if the sides of the ladder were J = 0 exchanges, *i.e.* π 's or σ 's, then the density of rungs, representing produced clusters (or clans, in the language of multiparticle dynamics [17]) of hadrons should be

$$\frac{dN_{\sigma}}{d\eta} \sim 2 \quad ? \tag{21}$$

Analyses of the multiplicity distributions has suggested in fact that a reasonable value for the mean multiplicity of hadrons per clan is about 4, leading to a total

$$\left(\frac{dN}{d\eta}\right)_{\rm hadrons} \sim 8 ?? \tag{22}$$

This value is a little high but not unreasonable, so that the hypothesis of pion/ σ exchange is at least not disfavored by data. Actually the clans or clusters may well be the products of $q - \overline{q}$ systems of mass 1–2 GeV, since the Yukawa coupling of pions to quarks, as we have seen, is quite large. However, the above estimates are at best semiquantitative, since the derivation was made in the weak-coupling, low density limit.

In the case of the hard, perturbative QCD Pomeron, the mechanism is quite similar. The production cross section for n QCD gluon jets into the phase space has no power-law dependence on s and is proportional to $(\alpha_s/\pi)^n$. The jets drop randomly, more or less, into the phase space [18]. If the p_t scale of all the produced jets is demanded to be within, say, a factor 2 of a common value m, then even after radiative corrections, real and virtual, the size and shape of the differential cross section remains the same (for the same reasons that the radiative corrections to the $e^+e^- R$ do not have corrections of order $\alpha \log$). The energy dependence of the QCD Pomeron is calculated to be

$$\sigma \sim s^{(12\alpha_s \ell n 2)/\pi} \tag{23}$$

leading to an inferred gluon-jet density per unit rapidity of

$$\left(\frac{dN}{d\eta}\right)_{\text{gluon jets}} \sim \frac{12\alpha_s}{\pi} \ln 2 \ . \tag{24}$$

However there is no Regge-pole in this case; the reason has to do with the lack of a p_t cutoff, in either infrared or ultraviolet, in the theoretical calculations. (I think, however, that it may be a mistake to leave the p_t cutoff out.)

Now let us turn to the Feynman-Van Hove picture [19] of the Reggeon. Clearly from the point of view of Feynman diagrams the exchange of a ladder in the s-channel dynamics is related to a bound-state equation when the process is viewed in the tchannel. In the case of QCD, if the sides of the ladder are a quark-antiquark pair, then there are only discrete bound states in the spectrum, no continuum. So the Feynman-Van Hove idea is to just model the ladder as the exchange of the entire "tower" of resonances that are the presumed output of the ladder equation.

The members of a "tower" are those which satisfy the same radial equation after the angular-momentum partial-wave decomposition (in the *t*-channel) is made. The strength of the angular-momentum-barrier term $L(L + 1)/r^2$ can be continuously mapped (actually via analytic continuation) from one state to the next; hence their output wave functions can as well.

Table 2:

	I = 0	I = 1	$(mass^2)_{I=0}$	$(mass^2)_{I=1}$
$3S^1$	$\omega(783)$	ho(770)	$0.69 { m ~GeV}^2$	$0.59 { m ~GeV^2}$
$3P^2$	$f_2(1270)$	$a_2(1320)$	1.61	1.74
$3D^3$	$\omega_3(1670)$	$ ho_{3}(1690)$	2.79	2.86
$3F^4$	$f_4(2050)$		4.20	

The most important towers are the ρ and ω families. Their properties are shown in the following table:



Figure 7: The Chew-Frautschi plot of spin versus squared mass: Regge trajectories. The ρ and ω trajectories are shown, as well as the conjectured trajectory (dashed) for the soft Pomeron.

The spin and masses are related to each other as shown in Fig. 7. The interpolating equation

$$J = \alpha(M^2) \tag{25}$$

is called the Regge trajectory.

With these preliminaries, we now proceed to the Feynman-Van Hove construction. First choose s and t small and compute individually the contributions of the members to the scattering amplitude. They will have the form:

$$A_J \sim \left[\left(\frac{s}{s_0}\right)^J + \cdots \right] \frac{G^2(J)}{[t - M^2(J)]} .$$
(26)

The three dots denote the nonleading contributions to the expansion of the angulardistribution Legendre function

$$P_L(\cos\theta) = (\cos\theta)^L + \cdots .$$
(27)

The leading $(\cos \theta)^L$ becomes eventually an s^J contribution [†], and the neglected terms will be nonleading ("daughter") trajectories. Now when $s/s_0 \ll 1$ the series above can be expected to converge. But we are interested in very large s, and a continuation is needed. It is here that, provided the coefficients $G^2(J)$ and $M^2(J)$ are known as function of J and have appropriate behavior, the machinery of the Watson-Sommerfeld transform can be used.

To see how the Watson-Sommerfeld transform works, consider the series

$$f(x) = \sum_{n=0}^{\infty} (n+1) x^n$$
 (28)

and make believe that we do not know how to sum it (and that, even though we are so stupid, we do know about Watson and Sommerfeld). All we want from the series is how it behaves at very large x. Elementary, says our dear Watson. Just write f(x)as follows:

$$f(x) = \frac{i}{2} \int_C \frac{dn(n+1)(-x)^n}{\sin \pi n}$$
(29)

with the contour encircling the poles on the positive real axis, as shown in Fig. 8. For x < 1 open up the contour along the imaginary axis (Yes, this is legal). Then let x get large; the integral is still defined. Finally with x large push the contour to the left—the further the better—picking up residues as appropriate from poles along the negative real axis. The contribution furthest to the right gives the leading behavior for large x. Note that in our example this occurs at n = -2, because the numerator n+1 has a zero at the first natural locale at n = -1.

Now we do this with our Reggeon amplitude, with the n plane becoming the J plane. The same thing happens except for an extra singularity when

$$t = M^2(J) \tag{30}$$

which as we saw is usually inverted as

$$J = \alpha(t) \tag{31}$$

$$t - M^2(J) = \frac{dM^2}{dJ} (J - \alpha(t)) + \cdots$$
 (32)

[†]Note that (for equal masses m for the external particles) the t-channel $\cos \theta$ equals $(s-u)/(t-4m^2)$.

i.e.



Figure 8: Watson-Sommerfeld transforms: (a) Complex n-plane for our simple example; (b) complex J plane for the Reggeons.

Note that as t becomes timelike this amplitude exhibits poles due to the factor $\pi/\sin \pi \alpha(t)$.

$$\sum_{J} A_{J} \to \left(\frac{-s}{s_{0}}\right)^{\alpha(t)} \frac{\pi}{\sin \pi \alpha(t)} \left(\frac{dJ}{dM^{2}}\right) G^{2}(\alpha(t)) + \cdots$$
 (33)

So the bottom line is that the scattering amplitude has a pure power-like sdependence which depends upon t. As t increases from zero to more spacelike values, the s-dependence weakens. This leads to the prediction of a shrinking of the width of the diffraction peak, which is observed. Furthermore the rate of shrinkage is determined by the Regge trajectory (and vice versa). The physics of this is that the location of the n^{th} "rung" in the transverse impact space is only correlated with its nearest neighbor. This leads to a random-walk mechanism, with the squared distance in impact space growing linearly with the number of rungs, hence in proportion to log s. The slope (in t) of the Regge trajectory determines the step size (the mean distance Δb between adjacent clusters/clans in impact space. A simple calculation gives

$$\langle b^2 \rangle = \langle b_0^2 \rangle + (\Delta b)^2 \frac{dN}{d\eta} \ln s .$$
 (34)

The formula for the shrinkage is

$$\frac{d\sigma}{dt} \sim e^{-b^2|t|} \sim F(t) \, s^{-(\Delta b)^2 (dN/d\eta)} = f(t) \, s^{2\alpha(t)-2} \,. \tag{35}$$

So putting this together gives

$$2 \text{ GeV}^{-2} = 2 \frac{d\alpha}{dt} \cong (\Delta b)^2 \frac{dN}{d\eta} .$$
(36)

For the ρ trajectory, with quarks on the sides of the ladder, the lowest order energy dependence of the amplitude would correspond to a bare α of 1/2, or s^{-1} cross

section dependence. Instead, from the Chew-Frautschi plot (Fig. 7), the dependence is $s^{-1/2}$, leading to a density of rungs, or clusters in the lego plot, of about 1/2 per unit rapidity. This gives a step size of 0.4 fermi for the ρ -exchange random walk. For the QCD hard Pomeron the random-walk is supposed to be very small; it is in fact questionable whether there is any shrinking of the diffraction peak at all.

The most interesting case is that of the soft Pomeron, which controls the total cross section at high energies. From an s-channel point of view, soft Pomeron exchange is just a fancy way of saying that at high energies the soft strong interactions are really strong, so strong that low partial waves are almost completely absorbed. So there must be shadows cast—namely a considerable amount of elastic scattering. Were the hadrons black discs, the elastic and inelastic scattering would be equal in magnitude. In practice the elastic cross section is 20-25% of the total, but the ratio is rising with energy and the shape of the elastic scattering angular distribution is changing with energy in a way that looks more like black-disk [20].

Likewise, the shrinkage of the diffraction peak can be simply expressed in s-channel language as the natural growth with energy of the size of the black disk. Absorption depends on the momentum density of the partonic matter in the disk; if the momentum per square fermi carried by the projectiles is large compared to 1 GeV, then one can expect lots of absorption. Near the edge of the disk it is presumably of that order. But when the energy of the projectile is increased tenfold, that momentum density likewise increases tenfold, so the critical region defining the edge of the disc has to move outward [21].

So it is arguable that there is no need to invoke t-channel Regge pole ideas to describe the soft Pomeron. Nevertheless Donnachie and Landshoff have given rather persuasive arguments [22] that the soft Pomeron is in fact a Regge pole, not cut, with intercept at t = 0 of 1.08, and with a slope $d\alpha/dt = 0.25 \text{ GeV}^{-2}$. This object couples to the constituent quarks, and would naturally be connected with dynamics at about the GeV mass scale. If the soft Pomeron is a Regge pole, there should be (cf. Fig. 7) a spin 2 resonance when $\alpha(M^2) = 2$. This occurs at M at about 2 GeV. And indeed there is a candidate 2^{++} state at about 1900 GeV which seems to couple strongly to the soft Pomeron; it is seen [23] in $pp \to pp X(1900)$, with the X decaying to 4 pions. The width is large, about 350 GeV.

Landshoff and Nachtmann [24] models the soft Pomeron as a bound state of two "constituent" gluons. Therefore the 2⁺⁺ resonance on the soft-Pomeron Regge trajectory would naturally be identified as a gluon-gluon bound state (gluonium). But the large width to pions might pose a problem. It alternatively might be a bound state of pions/ σ 's/constituent quarks. But no matter what it is, the hypothesis that the soft Pomeron is a Regge pole is the main issue; it implies very special ladder-like, bound-state underlying dynamics.

5 Soft and Hard Diffraction

The basic definition of a diffractive process is that it should occur via the strong interactions and that there should exist a large "rapidity-gap", not exponentially suppressed, in the final state. By rapidity-gap is meant that within some interval $\Delta \eta$ in the lego plot, with $\Delta \eta \gg 1$, there are *no* hadrons produced. (This implies that the initial cms energy be large, typically greater than 30–100 GeV, in order to have enough phase-space to produce the gap.) By "not exponentially suppressed" is meant that the probability of finding a large rapidity gap is not a strongly decreasing function of $\Delta \eta$. For example the process $e^+e^- \rightarrow \pi + \pi$ exists at all energies, and at high energies it clearly has a big rapidity gap in the final state. But the cross section falls like a power of *s*, hence exponentially with rapidity separation of the pions. In general e^+e^- single-photon annihilation into hadrons is non-diffractive.

We shall distinguish further two kinds of diffractive processes—hard and soft. Hard diffraction contains jets in the final state; soft diffraction generically does not. Evidently hard diffraction will be naturally associated with exchange of the perturbative QCD Pomeron, and soft diffraction with the soft Pomeron. In general, "Pomeron exchange" is interchangeable with "diffraction", although there clearly are additional theoretical overtones when the former term is used.

The most prominent diffractive process and the one with the longest history is elastic scattering. Closely related to it is the total cross section, because it is just the imaginary part of the forward elastic amplitude. The next most important process is single diffraction-dissociation, which is rather well measured up through TeVatron collider energies [25]:

$$\sigma(p\overline{p} \to pX) + \sigma(p\overline{p} \to \overline{p}X) \approx 9.5 \pm 0.5 \ mb \ . \tag{37}$$

What happens is that one of the projectiles is excited into a massive final state while the other remains intact. The projectile excitation can be viewed in optical model terms—it is just the differential absorption across the impact plane of the constituents. Suppose it is the proton which is excited. The quarks within the proton which are near the collision axis get absorbed more than those further away; therefore the wave function of the outgoing proton system is no longer the ground state wave function and there must be excitation [26]. The particle distribution in the lego plot is as shown in Fig. 9. Double diffraction dissociation exists as well; however it has not yet been thoroughly studied at TeVatron collider energies.

As we have already indicated, these processes are often described in *t*-channel terms: "the exchange of a Pomeron". While an optical-model description emphasizes the feature of absorption, thereby making the Pomeron a rather shadowy object, it is nevertheless the case that in high-mass diffraction, the Pomeron carries across a large amount of energy and momentum in order to create the excitation. At least, there are a class of reference frames where this is true, as can be seen by comparing the



Figure 9: (a) Single diffraction dissociation, and (b) the lego plot for the final state particle distribution.

QED situation of photon-exchange, say in an electron proton collision, with the QCD situation of Pomeron exchange. The shadowy view of the diffraction is analogous to looking at the *e-p* collision as Coulomb excitation. The electron is a left mover; all proton fragments are right movers, and predominantly transverse momentum is exchanged. In the hadron collision, the \overline{p} , say, is the left mover and is not excited. Instead it elastically scatters from a low-x subsystem Q of the right-moving proton. This innocuous elastic scattering can create in fact a large final state mass. To see this, calculate $E - p_z$ for the excited system:

$$(E - p_z) = \frac{M^{*2}}{2E} \cong \frac{(p_t^2 + m_Q)^2}{2x_Q E} .$$
(38)

This can be very large when x_Q is very small.

The more dynamical view of the Pomeron is analogous to viewing the e-p process in the cms system of the excited proton; one then has a transverse photon colliding with the initial-state proton and producing hadrons. What is the analogous picture for the Pomeron? What is the analogue of the Weiszacker-Williams photon? Operationally it should be possible to describe the Pomeron in terms of the quanta carrying its large longitudinal momentum. In fact it is possible, for each process, to operationally define a parton distribution for the Pomeron and to expect these partons to generate hard collisions. This is the seminal suggestion of Ingelman and Schlein [27], which launched the present program of studying hard-diffraction processes.

However, despite the above argument, and its great importance in initiating the experimental study of the short-distance structure of diffractive processes, I am not now convinced that the notions of exchanged Pomeron and of Pomeron partondistributions are necessary to describe the phenomena, and that instead it is possible, and in fact more simple, to hold to the *s*-channel language. That is what I will try to do in what follows. In particular the picture of the elastic scattering of the subsystem as used above, needs to be examined in the "Weiszacker-Williams" class of reference frames.

In most of what follows we will focus on the new results on diffractive processes from HERA and the Fermilab collider. It will turn out that, not surprisingly, the simplest class of experiments which probe these issues are electron-proton collisions. The advantages for the $p\bar{p}$ collisions studied at CERN and at Fermilab lies in the higher energy scale, and we save them for the end.



Figure 10: (a) Diffractive electroproduction and (b) the lego plot for the final-state particle distribution.

A very clean way to define the Ingelman-Schlein "structure function of the Pomeron" is in the diffractive process shown in Fig. 10. The undissociated proton is tagged with a special "Roman pot" detector far downstream of the collision point and very near the beam. What is left is the collision of the electron (or virtual photon) with the exchanged Pomeron. The structure function of the Pomeron, however shadowy it really is, is operationally defined by this process. However, let us look at this process from the s-channel point of view. As discussed already in Section 3, we may for most collisions regard the virtual photon as Blike. It dissociates into a fast, passive quark-parton q (analogous to the \overline{b} in the B) and a slow, collinear constituent antiquark \overline{Q} (or vice versa; $q \leftrightarrow Q$ and $\overline{Q} \leftrightarrow \overline{q}$). On arrival at the collision point all that need happen in the diffractive process is that this constituent antiquark *elastically* scatter from the proton. "The Pomeron is exchanged", the rapidity gap is formed, and the cross-section is easy to calculate: the ratio of the rapidity-gap process to the total is just the ratio of elastic (constituent) \overline{Q} -proton scattering to total \overline{Q} -proton scattering. This latter ratio, probably of order 10%, give or take a factor 2, should be essentially independent of the scaling variables Q^2 and x—to the extent that the aligned-jet mechanism dominates the ordinary structure function F_2 , and that the total and elastic quark-proton cross sections do not vary with energy.

All this is quite consistent with the data, to the best of my knowledge. Clearly the description could hardly be simpler, and suggests that experimentally the determination of the *ratio* of diffractive cross section to nondiffractive cross section as function of x and Q^2 (and eventually t, the squared momentum transfer to the proton) is not only experimentally convenient, but also quite relevant theoretically. Also note that the elements of the argument involve from beginning to end concepts well beyond perturbative QCD: the constituent (anti)quark and its black-disk-like absorption on the proton.

A most interesting question remaining is how to describe this utterly simple result in the t-channel language, and to understand why this picture seems to be not inconsistent with a variety of explanations of the same phenomenon which use only the machinery of perturbative QCD.

So let us look again at the diagram for the process as described above. As drawn in Fig. 11a, we see the photon dissociation into $q\overline{Q}$ followed by the elastic scattering of \overline{Q} from proton, with the Pomeron momentum pure spacelike. In general the subprocess looks simply like

$\gamma + \text{Pomeron} \to q + \overline{Q}$ (39)

via a simple diagram. So let us interchange the kinematics of the gamma and Pomeron, and go to the frame (Fig. 11b) where the γ momentum is pure spacelike (a Coulomb photon). The time ordering is changed, so now it appears that the proton emits a virtual Pomeron, and then the Pomeron dissociates into a $Q - \overline{Q}$ pair (note that both Q and \overline{Q} from this point of view must be considered constituent quarks. This $Q - \overline{Q}$ system is anything but pointlike, because its mass is small; tis limited to at most a few GeV² because the beam proton is not dissociated. Then the electron finds in this mess a quark-parton and scatters from it at high Q^2 . The secondary quark-parton has high p_t and exits.

This is what the diagram says. But clearly the diagram is not very reliable. What is reliable? At best it is only the last two sentences in the previous paragraph. The



Figure 11: Two views of the diffractive electroproduction process: (a) Photon carries momentum; Pomeron exchanged; (b) Pomeron carries momentum; photon exchanged.

preceding words can be deleted with no loss of comprehension. But where then does the rapidity-gap come from? So let us try again, and start from the beginning. We find that in the above reference frame the electron first strikes the proton and knocks out the quark-parton q. As this q exits the proton it picks up a polarization-cloud with quantum-numbers of a \overline{Q} , because of confinement [28]. Does this cloud interact with the proton remnants or not? If it does not, the rapidity-gap is formed; if it does, an ordinary deep-inelastic final state is formed. The bottom line is that the Pomeron phenomenon in this frame appears to occur in the evolution of the *final* state, not the initial state as the (unreliable) diagram would suggest.

To see this a little better, look at the phase-space picture for the process, in the HERA laboratory frame of reference (Fig. 12). The phase-space location of the \overline{Q} "polarization cloud" is shown as well as the location of the struck-quark jet with $p_t = Q$. The separation of the \overline{Q} -cloud and the jet in the lego plot is by an amount of order log Q. The color separation occurs only between the jet and the \overline{Q} -cloud. If the frame of reference is chosen so that $\eta = 0$ is to the right of the \overline{Q} -cloud region, then the hadronization in the neighborhood of the \overline{Q} -cloud occurs earlier than the creation of the rapidity-gap, because the momentum scale for the quanta which would fill the candidate gap region is larger in this frame than for the \overline{Q} -cloud quanta, and because the time scale for evolution of the final-state system is in proportion to the momentum scale. If the frame of reference is chosen so that $\eta = 0$ is to the region, we revert to the original, simple "aligned-jet", \overline{Q} -elastic-scattering description.

Thus far we have investigated rapidity-gaps in the HERA data for final states which do not contain jets (The struck-quark jet we discussed above is an artifact of kinematics and can be eliminated by changing the reference frame to a collinear



Figure 12: Lego plot for the diffractive final state showing the locations of $\eta = 0$ for the choices of reference frame made.

proton-virtual-photon frame.) The Pomeron we have considered is therefore the soft Pomeron. (In more general terms we have only discussed soft diffraction.) But hard diffraction can also be expected in the HERA data—and may have been observed already. The interpretation of hard diffraction most naturally does use gluon exchange. It should occur when the virtual-photon configuration at arrival is Υ -like, not *B*-like. As discussed in Section 3, the $q - \overline{q}$ pair have in that case comparable longitudinal momenta and have p_t of order Q, *i.e.* large. In the final state there should be seen a leading dijet, which carries almost all of the virtual-photon momentum.

The lowest order diagram for this process is reliably considered in perturbative QCD, and is shown in Fig. 13a. It is just like Bethe-Heitler pair production in QED, and is dubbed photon-gluon fusion in QCD. A necessary condition for obtaining a rapidity gap via this mechanism is that no color be exchanged between the virtual-photon system (the $q - \bar{q}$ pair) and the target system. This is accomplished via exchange of a second gluon of opposite color to the first (Fig. 13b). Since no color is exchanged and the virtual photon system only hadronizes long after it exits the target, we expect a rapidity gap to occur between the photon fragments and the target fragments. The lego-plot is shown in Fig. 14.

A crude estimate of the frequency of this gap process is found by estimating the ratio of the two-gluon exchange to one-gluon exchange to logarithmic accuracy. The result is [29]

$$P_{\rm gap} = \frac{\sigma(2 - \text{gluon})}{\sigma(1 - \text{gluon})} \approx (\text{const.}) \left(\alpha_s(t)\ell n t\right)^2 \approx \text{const.} \approx 0.1 .$$
(40)

As for the soft-diffraction, aligned-jet mechanism, this ratio should be essentially independent of Q^2 and x. It also should be generalized to exchange of the entire



Figure 13: Diagrams for (a) Hard dijet electroproduction, and (b) its hard-diffraction counterpart.

hard-Pomeron QCD gluon ladder. But there is an additional subtlety [30]. If the momentum transfer t is small compared to Q^2 there is suppression of the gap probability estimated above. This occurs because the one-gluon exchange itself is suppressed by an extra factor t/Q^2 due to the smallness of the color dipole moment of the $q\bar{q}$ relative to the distance scale $t^{-1/2}$ probed in the scattering process. The two-gluon exchange is then doubly suppressed, leading to the t/Q^2 factor appearing in the gap probability as well.

Hard diffraction has also been seen in proton-antiproton collisions at the CERN and Fermilab colliders. Indeed the discovery of hard diffraction occurred at the CERN collider in the UA8 experiment led by Schlein. However, we will begin with the Fermilab experiments, since they are in my mind simpler to interpret.

In both the CERN and Fermilab experiments, two coplanar hadron jets are detected (Fig. 15). And in both experiments a rapidity-gap is of course seen. But for the Fermilab experiments the rapidity gap is between the two jets, and for the UA8 experiment the gap is between a pair of right-moving jets and a left-moving undissociated Roman-pot proton.

The two-jet process seen at the TeVatron is, neglecting rapidity-gap issues, just the Rutherford-like scattering of a left-moving parton from a right-mover via single gluon exchange. But exchange of color will not allow a large gap to form. And just as discussed for the HERA case, exchange of a second gluon of opposite color can lead to a colorless exchange and the possibility of formation of a rapidity gap. The price paid for exchange of the second gluon is again of order

$$\frac{\sigma(2 - \text{gluon})}{\sigma(1 - \text{gluon})} \approx 0.1 \tag{41}$$



Figure 14: Lego plots of the final state particle distributions for the processes described in Fig. 13.

as in the previous case. However, there is a difference. Because both projectiles are complicated extended objects, any interaction of spectator partons from the left moving disk with spectators from the right-moving disk can lead to extra particle production and a filling in of the gap. If the partons are uncorrelated in the impact plane the price to be paid, called $\langle S^2 \rangle$, the "survival probability of the rapidity gap", is a factor of about 10. It is calculable from elastic scattering [29,31]; after a Fourier transform, the imaginary part of the elastic scattering amplitude measures the probability that at a given impact parameter the incident projectiles go through each other without interacting. Note that were the partons tightly clustered around constituent quarks, the survival probability might be even lower.

In any case the expectation is that the probability for finding a rapidity gap between jets separated in the lego plot by an amount $\Delta \eta$ should fall exponentially for small $\Delta \eta$, but then level off and be roughly constant, at a value which was conservatively set (ahead of the data [29]) as between 10^{-2} and 10^{-4} . The expectation is borne out by the data from both CDF [32] and D0 [33]. The level is about 1%, indicating not much if any impact plane correlation of the spectator partons.

Finally we consider briefly the first of the hard-diffraction measurements, the UA8 data on dijet production. From the t-channel point of view, the subprocess (Fig. 16a) is Pomeron + parton \rightarrow jet + jet, where the Pomeron has a t which is neither very large nor small: $t = 2.5 - 3.5 \text{ GeV}^2$. Ingelman-Schlein parton distributions for the Pomeron [27] are introduced, and then determine from the analysis of the kinematics of the two-body hard-collision subprocess. What is seen [34] is that most of the structure function of the Pomeron has a structure x(1 - x) anticipated (by Donnachie and Landshoff [35]) as due to the Pomeron coupling to quark-antiquark, in



Figure 15: Lego plots for (a) CERN UA8 final state, (b) TeVatron CDF and D0 final states, and (c) a slight generalization of each.

a way similar to what was alluded to in the HERA soft-diffraction process. However, in addition to this contribution UA8 finds an extra "superhard" component to the Pomeron structure function, consistent [‡] with a δ -function: $0.2 \delta(1-x)$. This Pomeron is no ordinary hadron! What is going on? From an s-channel point of view, one would be tempted to adapt the aligned-jet argument to this case. The picture might be as in Fig. 17. The high p_t dijet seen in the final state is made via splitting of an initial-state parton. But the coupling of the dijet to the Roman-Pot proton is via soft Pomeron as shown. One gluon splits collinearly into $q\overline{Q}$, with the \overline{Q} a constituent antiquark as before. The constituent antiquark elastically scatters from the proton as in the HERA example.

On the other hand, the delta-function component might be the hard Pomeron (two gluons) being exchanged. This could lead to the delta-function-like structure function, because one of the gluons (to logarithmic accuracy) carries a much larger share of the momentum than the other. (The QED analogue is two-photon exchange at large t. Exchange of the second photon is relatively "soft" and simply contributes a Coulomb phase to the amplitude; *i.e.* the second order amplitude is $(i\alpha \log)$ times the first order amplitude [29].)

[‡]To logarithmic accuracy, one may replace $\delta(1-x)$ with $(1-x)^{-1}$; which may be more reasonable.



Figure 16: Mechanisms for the CERN UA8 hard-diffraction process: (a) Partons in the Pomeron scatter from partons in the \overline{P} , and (b) a constituent quark in the \overline{p} elastically scatters from the proton in conjunction with dijet production.

In any case, the net effect is to have the color-singlet amplitude be proportional to the lowest-order, one-gluon-exchange amplitude, independent of the kinematics. This essentially means that the Pomeron behaves as if it is a single gluon, not two (or more).

The same effect can be anticipated for the Fermilab situation, slightly generalized. Were there two jets seen on one side of the gap and one on the other, the simple estimates made for the 2-jet TeVatron data should survive [37]: the ratio of the cross section for the diffractive final state to the nondiffractive cross section, for identical trijet kinematics, should again be about 1%, independent of the rapidities and transverse momenta of the jets (provided the momentum transfer is not small compared to the jet p_t 's). On the other hand the process could be reinterpreted again in the Ingelman-Schlein way as a hard process involving constituents of the Pomeron. But again one may see from the preceding argument that the Pomeron acts as if it were a single gluon as far as the kinematic dependencies are concerned, *i.e.* it has a δ -function component to its structure function.

Note however that in all the applications to hadron-hadron collisions, one cannot expect that the parton description of the Pomeron structure is simple, because the factor $\langle S^2 \rangle$, the survival-probability of the rapidity gap, must enter that description. This introduces a factor-10 renormalization of the Pomeron structure function and in addition introduces the physics of spectator interactions into the description of Pomeron structure.



Figure 17: Mechanism for the superhard-Pomeron final state seen in the UA8 data.

6 Concluding comments

There are no real conclusions to be made here. These lectures have been too brief to draw firm conclusions. But the main points I would emphasize are the following:

- 1. The constituent-quark structure seen in spectroscopy needs to be better understood. An attractive hypothesis is that this structure is connected with the spontaneous breaking of the strong chiral symmetry, and leads to the observed constituent quark masses and a large Yukawa coupling to pions. The behavior of the Gottfried sum rule and spin sum rules is also suggestive of the relevance of a "pion cloud" surrounding constituent quarks.
- 2. The constituent quark structure is evidenced in high energy soft collision processes via the additive quark model. In addition, the soft Pomeron seems to couple to constituent quarks. This is a slightly stronger phrasing of the essence of the additive quark model.
- 3. Diffractive processes can shed light on some of these issues. Soft diffraction, which seems to occur even in HERA deep-inelastic processes, may well probe the nature of constituent quarks. Hard diffraction is more amenable to perturbative-QCD treatment. But the two kinds of diffraction eventually must mix, and it is not clear where the dividing line is—or will be. Low- Q^2 HERA data is sure to be important, because experimentally real photoproduction looks like soft-Pomeron hadron physics.
- 4. There was unfortunately no time left in these lectures to mention disoriented chiral condensate (DCC). DCC is strong-interaction vacuum with a chiral order parameter which is not oriented in the vacuum direction. It has been conjectured [36] that it can be produced in high energy hadron-hadron collisions or ion-ion collisions. Its decay is into coherent, semiclassical pulses of pion field

of definite (cartesian) isospin. This in turn leads to anomalously large fluctuations in the charged-to-neutral ratio of produced pions. If DCC exists it will be coupled to and/or associated with constituent quarks, and is therefore also relevant to everything we have discussed. The experimental search for DCC is at present my principal occupation, and a prime reason for the sketchiness of these lectures.

Clearly there are plenty of fundamental questions left to answer. It is gratifying that there is a growing experimental interest—and capability—to help provide the answers.

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References

- 1. M. Goldberger and S. Treiman, Phys. Rev. 110, 1178 (1958). For a derivation based on chiral symmetry breaking, see *e.g.* J. Bjorken and S. Drell, "Relativistic Quantum Mechanics" (McGraw-Hill, New York, 1964), Chapter 10.
- 2. A. Manohar and H. Georgi, Nucl. Phys. **B234**, 189 (1984). See also A. Manohar, these proceedings.
- 3. There are useful discussions in Field's book: R. Field, "Applications of Perturbative QCD," Addison-Wesley, Menlo Park, California (1989).
- 4. This "aligned jet" picture has slowly evolved over a long period; cf. e.g. J. Bjorken, AIP Conference Proceedings No. 6, Particles and Field Subseries No. 2, ed. M. Bander, G. Shaw, and D. Wong (AIP, New York, 1972) for something early, and L. Frankfurt, Proceedings of the International Workshop on Deep Inelastic Scattering and Related Subjects," Eilat, Israel (February 1994), ed. A. Levy (World Scientific, Singapore), p. 354, for something more recent.
- 5. For a recent review, see F. Tessarotto, Proceedings of the 24th International Symposium on Multiparticle Dynamics", Vietri sul Mare, Italy, September 1994, ed. A. Giovannini, S. Lupia, and R. Ugoccioni (World Scientific, 1995, Singapore), p. 567.
- 6. J. Bjorken, Acta Physica Polonica B23, 56 (1992).
- 7. T. Akesson et al., Z. Phys. C34, 163 (1987).
- 8. J. Alitti et al., Phys. Lett. **B268**, 145 (1991).
- 9. F. Abe et al., Phys. Lett. **D47**, 4857 (1993).
- 10. See e.g. T. Sjostrand and M. Van Zijl, Phys. Rev. D36, 2019 (1987) and references therein.

- 11. M. Fatyga, whom I thank for most enlightening discussions.
- 12. F. Low, Phys. Rev. D12, 163 (1975); S. Nussinov, Phys. Rev. Lett. 35, 1286 (1975).
- 13. P. Collins and E. Squires, "An Introduction to Regge Theory and High Energy Physics," Cambridge University Press, Cambridge (1977).
- 14. R. Newton, "Scattering Theory of Waves and Particles," (Springer, 1982).
- 15. S. Frautschi, M. Gell-Mann, and F. Zachariasen, Phys. Rev. **126**, 2204 (1962).
- 16. D. Amati, S. Fubini, and A. Stanghellini, Phys. Lett. 1, 29 (1962).
- 17. A. Giovannini and L. Van Hove, Z. Phys. C30, 391 (1986). S. Lupia and R. Ugoccioni, Ref. 5, p. 458, and references therein, for an up-to-date discussion.
- 18. V. Del Duca, Phys. Rev. **D48**, 5133 (1993).
- 19. L. Van Hove, Phys. Lett. **24B**, 183 (1967).
- 20. M. Block, F. Halzen, and B. Margolis, Phys. Lett. B252, 481 (1990).
- 21. See for example, J. Bjorken, Ref. 6.
- 22. A. Donnachie and P. Landshoff, Nucl. Phys. B267, 640 (1986).
- T. Armstrong, et al., Phys. Lett. B228, 536 (1989); S. Abatzis et al., Phys. Lett. B324, 509 (1994).
- 24. P. Landshoff and O. Nachtmann, Z. Phys. C35, 405 (1987).
- 25. F. Abe et al., CDF Collaboration, Phys. Rev. D50, 5535 (1994).
- 26. M. Good and W. Walker, Phys. Rev. 120, 1857 (1960).
- 27. G. Ingelman and P. Schlein, Phys. Lett. B152, 256 (1985).
- 28. For a description of this kind of space-time evolution, see J. Bjorken, Lecture Notes in Physics 56' Current Induced Reactions, ed. J. Koerner, G. Kramer, and D. Schildknect, Springer-Verlag (N.Y.), 1975.
- 29. See J. Bjorken, Phys. Rev. D47, 101 (1993).
- 30. J. Collins, L. Frankfurt and M. Strikman, Phys. Lett. B307, 161 (1993).
- 31. E. Gotsman, E. Levin and U. Maor, Phys. Lett. **B309**, 199 (1993).
- 32. F. Abe et al., Phys. Rev. Lett. 74, 855 (1995).
- 33. S. Abachi et al., Phys. Rev. Lett. 72, 2332 (1994).
- 34. A. Brandt et al., Phys. Lett. **B297**, 417 (1992).
- 35. A. Donnachie and P. Landshoff, Nucl. Phys. B303, 634 (1988).
- 36. For recent discussions, see J. Bjorken, Proceedings of the Workshop on Continuous Advances in QCD, Minneapolis, MN, February 1994, ed. A. Smilga (1994); also A. Krzywicki, Proceedings of the 29th Rencontre de Moriond, Les Arcs, France, March 1994, ed. J. Tran Thanh Van (Editions Frontieres, Singapore).
- 37. J. Bjorken, Eilat Workshop (Ref. 4), p. 151.