Recent Progress in QCD

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The past few years have seen remarkable progress in the theory and phenomenology of QCD, bringing perturbative and nonperturbative methods into closer contact with each other and with experiment.

I. INTRODUCTION: QCD A GROUP PORTRAIT

In this talk, I will summarize some recent developments in QCD, concentrating on, but not limited to, topics discussed at the this meeting. Details, of course, can be found in the talks themselves, presented in lively sessions organized by Lance Dixon and Joey Huston for perturbative QCD, and by Paul Mackenzie and Claude Bernard for nonperturbative QCD. The future of the field is in the already-advanced convergence of these topics, sometimes thought of as nearly independent. Progress on heavy quark physics is described in the proceedings from separate parallel sessions and plenary talks. I must necessarily pass over some of the most interesting recent advances in other closely related fields as well for lack of space.

I have tried in the following to interleave perturbative and nonperturbative treatments of QCD dynamics. Let me begin with a few general comments on the place of QCD studies in high energy physics.

Why QCD? By now it is clear that QCD is a "correct", or phenomenologically relevant theory, at least the way classical Maxwell theory is "correct". Classical electromagnetism is an effective theory appropriate to the limit of many photons; quantum chromodynamics might be the long-distance limit of some more fundamental, underlying dynamics. Its self-consistency, however, leaves us to free to study QCD in its own terms.

This is a fascinating, and dauting, task, despite the fact that QCD is defined through a single dimensionless parameter, α_s. The overall dimensional scale is determined by comparison with other interactions, for example, by measuring the strong coupling at the mass of the Z: $\alpha_s(M_{\rm Z})$. Its intrinsic interest aside, QCD has important "practical" applications, in the calculation of backgrounds to new physics and, for hadron colliders particularly, in predictions for new particle cross sections. From a theoretical point of view, however, what makes QCD so attractive is that it is a quantum field theory that requires all orders in perturbation theory and nonperturbative analysis to confront data at available energies.

QCD is the exemplary quantum field theory. QCD exhibits all of the classic quantum field-theoretic phenomena discovered in the sixties and seventies, including asymptotic freedom, confinement, chiral symmetry breaking and instantons. The problems of strong interactions that gave rise to QCD were also, in the same time period, the original inspiration for concepts of duality and strings. As we saw at this conference, the strong interactions, now understood as QCD, are once more the meeting ground for field theory, duality and string theory.

Again, QCD is evidently the correct theory of the strong interactions. Given its depth, however, "tests" of QCD should be thought of as tests, or perhaps better, "explorations", of quantum field theory itself. Complementary to the extraordinary accuracy of selected perturbative predictions in quantum electrodynamics are the broad predictions of QCD, interweaving nonperturbative and perturbative scales and phenomena.

II. QCD AT THE SHORTEST DISTANCES

It was the asymptotic freedom of QCD that first drew attention to gauge field theory as the unique description of the strong interactions at short distances. This theme continues to develop in current experiments, and to provide a basis for extrapolations between energy scales and the detection of signals for new physics. Let us begin with a run-through of the underlying methods [1].

A. Methods

Infrared safety. Infrared safe quantities are insensitive to long-distance effects, and may be calculated in perturbation theory. An infrared safe quantity may or may not be directly observable. The classic examples include the total cross section for e^+e^- annihilation to hadrons, and jet and event shape cross sections in e^+e^-. These can be written in the general form,

$$Q^2\,\hat{\sigma}_{\rm phys}(Q^2) = \sum_n c_n(Q^2/\mu^2)\,{\alpha_s}^n(\mu)\,, \tag{1}$$

with the c_n dimensionless functions of the ratio of the hard scale to the renormalization scale μ.

Factorization. Cross sections for deep-inelastic scattering (DIS) and for jet or heavy quark production in hadron-hadron scattering, are not purely perturbative, but appear as convolutions in parton momentum fractions of distributions $f_{a/h}(\xi,\mu)$ of partons a in hadrons h, with perturbative hard-scattering functions $\hat{\sigma}_{\rm PT}$,

$$Q^2\sigma_{\rm phys}(Q,x) = \sum_{\text{partons } a} \hat{\sigma}^{\rm PT}_a(Q/\mu,\alpha_s(\mu))\otimes f_{a/h}(\mu) = \sum_{\text{partons } a}\int_x^1 d\xi\;\hat{\sigma}^{\rm PT}_a(x/\xi,Q/\mu,\alpha_s(\mu))f_{a/h}(\xi,\mu)\,, \tag{2}$$

where in the second equality we have exhibited the convolution appropriate to deeply inelastic scattering (with $x=Q^2/2p\cdot q$). Corrections to Eq. (2) are suppressed by $\mathcal{O}(1/Q^2)$ [1]. In this formula, there is a complementarity between the roles of parton distributions and the hard scattering. The distributions $f_{a/h}$ are universal among hard-scattering processes, but particular to hadron h, while the functions $\hat{\sigma}^{\rm PT}$ are particular to the process, but universal among external hadrons. This last feature enables us to calculate realistic $\hat{\sigma}^{\rm PT}$ in "unrealistic", but technically manageable (infrared regulated) scattering processes, in which the initial state hadrons are partons.

Evolution. The physical cross sections of Eqs. (1) and (2) above must both be independent of the scales μ that define the parton distributions: $\mu d\sigma_{\rm phys}/d\mu = 0$. This self-consistency requirement is readily translated into the "DGLAP" equation for the evolution of parton densities,

$$\mu\frac{df_{a/h}(\xi,\mu)}{d\mu} = \sum_b P_{ab}(\xi/\eta,\alpha_s(\mu))\otimes f_{b/h}(\eta,\mu)\,. \tag{3}$$

Here, the dimensionless kernel P depends only on variables that are in common between the hard scattering functions and the parton distributions: α_s and the momentum fractions. The scale-independence of physical quantities and their relations can be studied systematically [3].

The idealized pattern for determining and applying the distributions may be summarized as follows. We measure one cross section, $\sigma_{\rm phys}$, at momentum transfer Q_0. Given a "next-to-leading" order calculation of its hard scattering functions $\hat{\sigma}^{\rm (NLO)}$, we determine NLO parton distributions $f^{\rm (NLO)}_{a/h}$ at $\mu=Q_0$. Using evolution, these can be used to predict $\sigma_{\rm phys}$ for any hard process at all Q.

The coefficients c_n in Eq. (1) are known for many processes to NLO [2]. They have been determined at NNLO for inclusive DIS and Drell-Yan [4], and even to three loops [5] for selected quantities. Generally, two loop corrections are available only for single-scale processes, and the calculation of two-loop corrections for genuine scattering diagrams is an as-yet unsolved, but actively studied, problem in QCD [6,7].

Perturbative QCD is most successful for inclusive processes, and/or single-scale semi-inclusive. Evolution in DIS is perhaps still the best illustration (see below). The current experimental situation in hadronic hard-scattering is

[1] I have described some of the technical background in Ref [1].

reviewed in Ref. [8]. In multiscale problems, logarithms of ratios of distinct but perturbative scales often require resummation to all orders. Formally beyond perturbative resummation, but not always less important numerically, are corrections suppressed by powers of the hard scale(s). In DIS, and a few other cases, these corrections can be desccribed by the operator product expansion. We shall encounter below "generalizations" of this famous technique, usually in the form of effective field theories.

B. Prime Examples

Tevatron Jets. The most impressive success in orders of magnitude continues to be the Tevatron inclusive single-jet and dijet cross sections [8–11], as illustrated in Fig. 1, which shows a plot from Ref. [12]. According to Eq. (2), these cross sections are of the form

$$\sigma_{p\bar{p}\to J+X} = \sum_{ab} f_{a/p} \otimes \hat{\sigma}^{(\mathrm{NLO})}_{ab\to J+X} \otimes f_{b/\bar{p}}\,. \tag{4}$$

We find a consistency with NLO theory and experiment at a few tens of percent, well within the overall systematic errors, over a range in which the cross section decreases by seven or so factors of ten. As the figure shows, reasonable choices of parton distributions (in this case CTEQ5HJ) can account even for the highest-momentum data, although a slight difference remains between the D0 and CDF data at the high end.

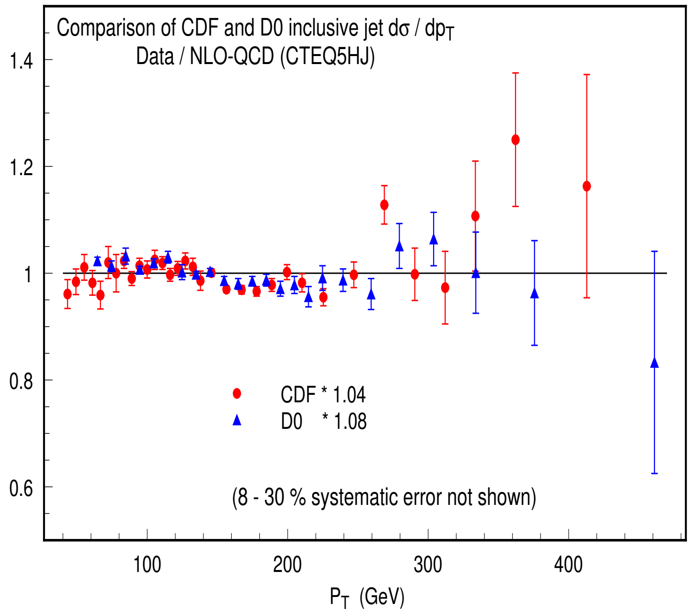

FIG. 1. A comparison of single-jet inclusive cross sections to NLO thoery from Ref. [12].

DIS Scaling violations. Next, we should cite measurements of DIS structure functions [13–15] for $\ell^\pm N \to \ell^\pm N$, through the cross sections

$$\frac{d\sigma^\pm}{dx dQ^2} = \frac{2\pi\alpha^2}{xQ^4}\,\left[\left(1+(1-y)^2\right) F_2(x,Q^2) - y^2 F_L(x,Q^2) \mp \left(1-(1-y)^2\right) F_3(x,Q^2)\right]\,, \tag{5}$$

at HERA, the Tevatron, SLAC and elsewhere, with $x = Q^2/2p\cdot q \quad q^2 = -Q^2 \quad y = Q^2/xs$. Because the evolution is universal in the factorized forms,

$$F_h(Q^2) = \sum_a C_a^{(\mathrm{NLO})}(Q^2/\mu^2) \otimes f_{a/h}(\mu^2)\ |_{Q=\mu}$$
$$\mu \frac{df_{a/h}}{d\mu} = \sum_b P_{ab}(\alpha_s(\mu)) \otimes f_{b/h}\ , \tag{6}$$

we may think of evolution as in, but not of, the nucleon, except perhaps in the smallest-x region, where target-dependent shadowing, in which partons begin to interfere with each other, comes into play.

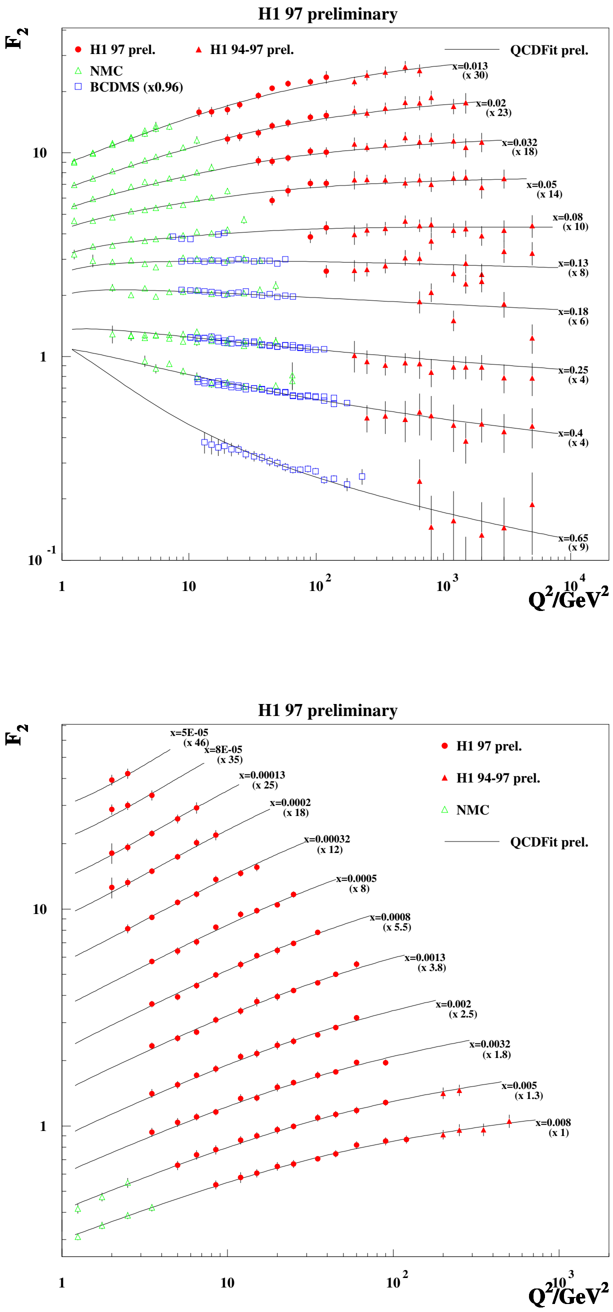

FIG. 2. DIS data from H1 (preliminary) [16], NMC and BCDMS.

The quality of the data, and of QCD fits to it, is illustrated by recent data on $F_2(x, Q^2)$ from H1 [16], shown in Fig. 2 for various fixed x as a function of Q^2. Notice in particular the sharp rise of F_2 with Q^2 for the smaller values of x. This is a prediction of the evolution through parton branching described Eq. (6). Deviations from such "DGLAP" evolution have been surprisingly difficult to find. The interesting results of a next-to-next to leading order analysis of DIS has been described in Ref. [17]. At the other end of the spectrum, the excess of events reported two years ago at the largest x and Q^2 has nearly disappeared [18].

Multijets, and event shapes. In e^+e^- annihilation, the Born level process $e^+e^- \rightarrow q\bar{q}$ produces two jets, as is the case as well for hadron colliders, while in DIS, the Born scattering produces a single jet. In all three "canonical" scattering processes, the cross sections for the minimal reactions are known at NLO. Beyond, this, NLO cross sections are available for two jets in DIS [19], for four in e^+e^- [20], and NLO three-jet cross sections are just now becomming available for hadron-hadron scattering [21]. Recent years have seen an explosion of data on these processes, and new results were discussed at this conference in [9–11,19,22], along with jet fragmentation properties [22–27]. The production of vector bosons associated with jets [28] is another important source of information about short-distance dynamics, and a reanalysis of W+1 jet cross sections has brought them back into agreement with theory over the past year.

Jets are theoretical-phenomenological constructs [29]. Their value is not that they are an exact reflection of short-distance reactions, but that, if they are defined properly, they are related to them in a calculable fashion. Jets are generally defined in terms of energy flow into some angular region, or in terms of interative clustering schemes for the momenta of observed particles. Any (sufficiently smooth [30]) quantity that is insensitive to the emission of zero-momentum particles, or to the collinear branching of finite-energy massless particles, can be used to define an infrared-safe cross section. Event shapes in e^+e^- annihilation or DIS are chosen for this property. The best-known is the thrust, defined for individual events as the maximum fractional projection of the momenta of observed particles on an axis, as that axis is varied about the unit sphere. Although event shape cross sections are infrared safe, they generally receive nonperturbative corrections that decrease rather slowly with energy, typically as a single power [31],

$$\sigma_{\rm phys} = \hat{\sigma}\,(1 + \mathcal{O}(1/Q)) \ . \tag{7}$$

The theory of these power corrections is a passageway between the short-distance perturbative, and the long-distance nonperturbative aspects of QCD. We shall have more to say about them below.

The α_s lineup. Any of the short-distance cross sections above allows a determination of the strong coupling. For example, in a factorized cross section, Eq. (2), given parton distributions f, we compare $\hat{\sigma}(\alpha_s)$ to experiment, and solve for α_s, typically evaluated at a renormalization scale equal to the factorization scale. Other, more refined, choices are related by perturbative corrections. For cross sections that are infrared safe, as jet or event shape cross sections at LEP [32], the comparison is even more direct, although power corrections, as in Eq. (7) must be taken into account for any precise determination [22,33]. Yet another way of determining α_s is from lattice QCD calculations of energy level differences in heavy quark systems. These may be related, on the one hand to the strong coupling, and on the other to experiment, which then determines the size of the coupling at a scale that grows with the heavy quark mass.

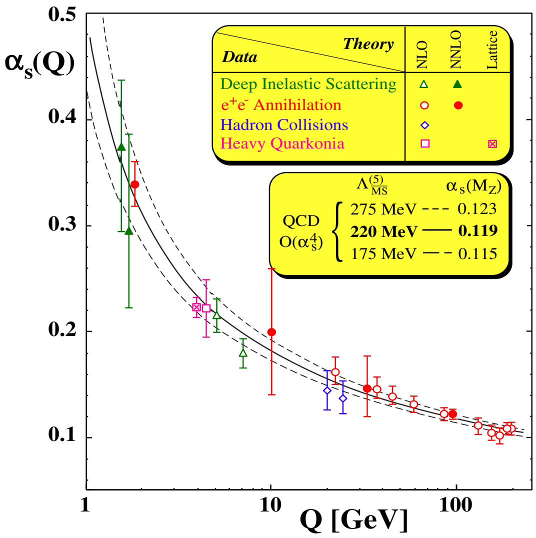

FIG. 3. The strong coupling measured at a different length scales, from Ref. [34].

For ease of comparison, the couplings can be evolved evolved to the mass of the Z, at which a "world average" of $\alpha_s(M_Z) = 0.119 \pm 0.004$ has been quoted in Ref. [34]. Since varied experiments measure the coupling over a wide range of scales, Fig. 3, which shows the variation of measured $\alpha_s(\mu)$ with μ illustrates one of the great successes of renormalized field theory.

C. Resummation

In the presence of multiple scales, high order contributions to hard scattering functions $\hat{\sigma}^{\rm PT}$ can become important. For example, in Eq. (2), the limit $x/\xi \rightarrow 1$ is associated with integrable singularities that usually enhance the cross section. These "threshold" singularities were resummed long ago for the Drell-Yan cross section, and have been applied for some time to top production at leading-log accuracy [47]. In QCD hard-scatterings, such as heavy quark and jet production, color exchange makes the resummation somewhat more complex beyond leading log. This problem has now been solved, and threshold resummation is understood in principle for a wide variety of hard scattering cross sections [48–50]. As a practical matter, differing approaches to the inversion of certain Mellin transforms can lead to differing numerical predictions. From a broad perspective, however, the main lesson is that the theory stays relatively close to NLO for cross sections like top production, even in the presence of superficially large corrections at higher order. Another important observation [47], made particularly clear in the very recent publication [50], is a marked decrease in factorization-scale dependence for resummed cross sections.

Another example, currently being discussed widely, is related to the data of Fig. 4, single-photon and pion inclusive cross sections measured by the fixed-target Fermilab experiment E706. So far, this data can be fit only by supplementing NLO with the old method of k_T smearing for the initial partons [51]. At the same time, it has been noted that the full range of fixed-target direct photon data may not be consistent among themselves [52]. It should also be noted that at collider energies, experiment and NLO agree, at least at the higher transverse momenta [35].

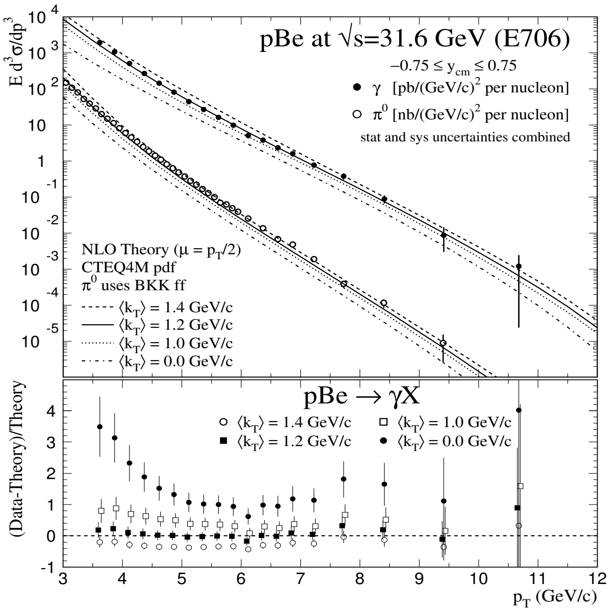

FIG. 4. E706 data versus NLO predictions and k_T-smearing.

Can the resummation of higher orders in QCD lead to an effective k_T smearing? For W or Z production at low transverse momentum, the answer is yes, and the formalism has existed for some time. In this case, fitting the cross section [36] requires the introduction of nonperturbative parameters that are accessible to experiment. In the past year or two, there has been some discussion on the best way of going about this [37,38], but the underlying theory is relatively well-understood. The same cannot be said for W or photon cross sections at high transverse momenta, because the logarithms whose resummation requires the nonperturbative input at $Q_T \sim 0$ cancel order-by-order in perturbation theory in the calculation of $\hat{\sigma}$ at high Q_T. Interesting observations on the relationship between the two regimes have been made, however [39], and futher progress in this direction can be anticipated in the coming year.

One of the major challenges in perturbative QCD is the development of a theory of these and related higher-order effects, and a method for estimating their importance. We shall return to developments on resummation below.

III. THE LONG AND SHORT OF HADRON STRUCTURE

A. Parton Distributions 1999

The parton distribution functions (PDFs) in Eq. (2) summarize the structure of hadrons as seen at short distances, one parton at a time. As the hadron is probed at ever shorter distances, the distributions evolve perturbatively according to Eq. (3), but the boundary condition for this evolution stands as a truely nonperturbative reflection of the dynamics that holds the hadron together. These parton distributions have been the subject of intense study since the recognition of approximate scaling in DIS structure functions thirty years ago.

Over the past decade, two groups, “CTEQ” and “MRS” whose memberships have themselves evolved somewhat, have undertaken coordinated “global” approaches to the determination of parton distributions, taking into account data from a variety of processes and momentum scales. The past year has seen the development of the latest, best fits of these two groups, MRST [40] and CTEQ5 [12], which were discussed and compared at the conference in [41].

Global fits test the self-consistency of factorized cross sections in the sense that the fits are overconstrained, and because they can be checked against experiments not incorporated into the fits. Nevertheless, the fits must be improved as the data improves, and as it extends to extreme values of fractional momenta. Surprises can occur, especially in regions where a particular parton distribution does not contribute at leading order. Examples of such refinements from 1998 involved the ratio of d to u quarks in the proton, as tested by the W asymmetry [42] and Drell-Yan [43]. Generally speaking, indirect constraints on parton distribution functions are tentative.

Other cases where the results of global fits have been rethought involve situations where there may be important higher-order or power ("higher-twist") corrections to the cross section. The extraction of PDFs requires a parameterization of such effects, assuming that they can be brought under control. Examples include the role of higher twist in the extraction of neutron PDFs from deuterium data [44].

Up to this cycle of global fits, the primary processes employed were DIS, Drell-Yan and direct photons, the latter thought to be especially valuable for constraining the gluon distribution, which does not appear at leading order in the other two. The data of E706 [45,46], however, a sample of which was shown above in Fig. 4, has thrown this neat picture into disarray, as it disagrees decisively with the predictions of Eq. (2) at NLO. A k_T smearing approach has been used in the MRST fits of last year [40]. The CTEQ5 fits abandon direct photons in favor of jet cross sections to constrain the gluon at large x [12].

Another interesting problem is the treatment of heavy quarks. For energies near a heavy-quark mass, m_Q, their production may be calculated directly as a hard process, part of $\hat{\sigma}$ in Eq. (2). For energies much above m_Q, however, it may be advantageous to treat the heavy quark as a parton, thus automatically resumming logs of m_Q^2/s. A number of schemes to make this transition have been proposed [12,40,53,54], to treat charmed quark production at HERA.

The sophistication of these considerations of global analysis, and the need to make accurate predictions at high energy raises the question of how to estimate uncertainties in the distributions [55]. In part to explore these issues, new sets of distributions have been produced based on DIS data only [56], and methods have been introduced to quantify uncertainties systematically through statistical analysis [57].

B. Spin and Off-diagonal Distributions

The past few years has seen a rebirth of interest in the high-energy physics of spin, which, with advances in the technology of polarized beams and targets, has made possible the systematic study of spin at the parton level [13]. The DIS cross section for a nucleon with spin s may be represented in the standard form $d^2\sigma/d\Omega dE = (\alpha_{\rm EM}^2/2mQ^4)(E_e/E_e')L^{\mu\nu}W_{\mu\nu}$, in terms of spin structure functions defined as

$$W_{\mu\nu} = W_{\mu\nu}^{\rm unpol} + \frac{i}{E_e - E_e'}\epsilon_{\mu\nu\lambda\sigma}q^\lambda s^\sigma\, g_1(x,Q) + \frac{i}{(E_e - E_e')^2}\epsilon_{\mu\nu\lambda\sigma}q^\lambda \left[p\cdot q s^\sigma - s\cdot q p^\sigma\right] g_2(x,Q)\,, \qquad (8)$$

The function g_1 has a particularly transparent interpretation in terms of quark helicity, $g_1(x,Q) = \frac{1}{2}\sum_f e_f^2 \Delta q_f(x,Q) + O(\alpha_s)$, where Δq_f is the difference between the distributions of quarks with helicity parallel to the hadron's helicity and against it. At this conference, precision data for g_1^p and g_1^n were presented by E155 [58].

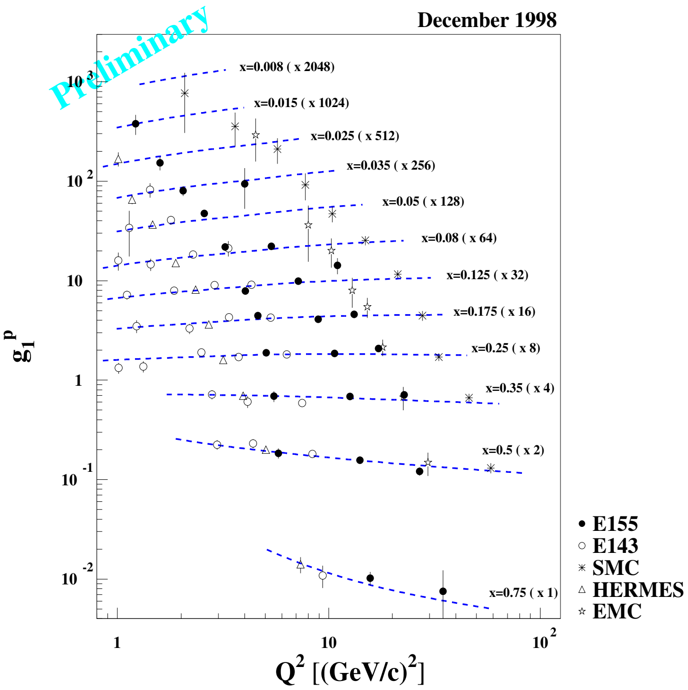

FIG. 5. E155 measurement of g_1 [58].

One result of these measurements is a test of of the benchmark Bjorken sum rule,

$$\int_0^1 (g_1^p(x) - g_1^n(x))\, dx = \frac{1}{6} \left| \frac{g_A}{g_V} \right| \left(1 + \sum c_n \alpha_s^n \right) , \tag{9}$$

which is now verified to a new level of accuracy.

The "spin" distributions of the nucleon, such as Δq_f, do not necessarily describe its full spin content, and the possibility of orbital angular momentum must also be taken into account. At the same time, for gluons the distinction between these two types of angular momentum is not gauge invariant. This problem notwithstanding, an attractive formalism for the description of orbital angular momentum has been proposed [59], in terms of form factors, $J_{q,g} = \frac{1}{2}\ [A_{q,g}(0) + B_{q,g}(0)]$, which arise in matrix elements of the energy-momentum tensor,

$$\langle p+\Delta |\, T^{\mu\nu}_{q,g}\, |p\rangle = \bar{u}(p+\Delta) \Big[\, A_{q,g}(\Delta^2)\, \gamma^\mu \bar{p}^\nu + B_{q,g}(\Delta^2)\, \bar{p}^\mu \frac{i\sigma^{\nu\alpha}}{2M} \Delta_\alpha - (\mu \leftrightarrow \nu) \,\Big] u(p)\, , \tag{10}$$

with $\bar{p} = p + \Delta/2$. At zero momentum transfer, the $J_{q,g}$ become expectation values of the angular momentum operators

$$\mathbf{J}_q = \int d^3x\, \left[\psi^\dagger \frac{\mathbf{\Sigma}}{2} \psi + \psi^\dagger \mathbf{x} \times (-i\mathbf{D}) \psi \right]\, , \quad \mathbf{J}_g = \int d^3x\, \mathbf{x} \times (\mathbf{E} \times \mathbf{B})\, . \tag{11}$$

The measurement of these off-diagonal matrix elements, unlike the diagonal matrix elements that define classic PDFs (see Eq. (16) below), require the measurement of exclusive, or semi-exclusive amplitudes, such as "deeply-virtual Compton scattering", $\gamma^*(Q^2) + N(p) \rightarrow \gamma + N(p')$, with $p^2 = p'^2 = m_N^2$. The measurement of such amplitudes is a challenge, but one that is of great interest for the Jefferson Laboratory facility, and there is a correspondingly vigorous theoretical program to study the factorization and evolution properties of off-diagonal distributions [60–62]. Off-diagonal matrix elements interpolate, in some ways, between inclusive and exclusive processes, and between parton distributions and hadronic light-cone wave functions [63].

The consideration of orbital angular momentum leads us to the doorstep of the long-distance, low-energy properties of hadrons, where progress has continued in lattice QCD and instanton studies.

C. Lattice Hadron spectra, quark masses

Lattice methods [64] approach QCD from a limit complementary to perturbation theory, and make possible direct calculation of long-distance properties of hadrons. Typically, lattice computations involve the evaluation of expectations of nonlocal products of operators, such as

$$C_J = \langle 0\ | J(x)\ J(0) |\ 0 \rangle\,, \quad J = \prod_i \bar{q}_i(x+\Delta)\, O_{ij}\ q_j(x-\Delta)\,, \tag{12}$$

with the q_i quark fields, where O_{ij} projects a set of quantum numbers. For $x \rightarrow \infty$ and $\Delta = 0$, the x-dependence is dominated by the energy of the lowest-lying state(s) of the relevant quantum numbers. As noted above, such studies can be used to set the scale for α_s by studying the hyperfine splitting in heavy quarkonia.

The numerical evaluations of expectations like (12) are said to be quenched, partially quenched, or fully unquenched, depending on how many species light fermions are allowed to fluctuate out of the vacuum. At the conference, progress was reported in quenched and unquenched lattice QCD, on the calculation of realistic spectra for hadrons, and of matrix elements of hadrons involving both heavy and light quarks [65,66]. Variations on this theme are now making possible the calculation of realistic decay matrix elements [66–68]. In addition, they allow the exploration of ideas on the mechanisms of confinement [69]. In alternative formulations, lattice methods may be applied to light-cone formulations of hadronic structure [63,70].

The current sophistication of lattice techniques, and the growing power of new machines, including Teraflop computers at the U. of Tsukuba and at Brookhaven(RIKEN)-Columbia, are making possible a new generation of investigations of chiral symmetry breaking.

Roughly speaking, chiral symmetry in QCD (also discussed at the conference in [71]) reflects the observation that gluon emission doesn't change helicities. As a result, in the absence of quark masses in the propagator, left- and right-handed quarks decouple altogether in perturbation theory. This means that a massless quark stays massless in perturbation theory. Chiral symmetry, however, is broken nonperturbatively at zero temperature, but restored at finite termperatures, a transition which can also be studied on the lattice [72,73]. At the same time, current algebra requires that the mass of the pion would vanish linearly with the lightest quark masses,

$$m_\pi^2 \sim m_{\rm light}\,. \tag{13}$$

Historically, this kind of relation has been difficult to realize on the lattice because of corrections that vanish as a low power of the lattice spacing in the continuum limit.

With the advent of ever more powerful machines, the method of "domain wall fermions" [74] makes possible new approaches to relations like Eq. (13). In the domain wall technique, the chiral symmetry is manifest, up to *exponential* corrections, at the price of introducing a 5th dimension, which we label by coordinate s, in the five-dimensional Dirac Lagrange density

$$\mathcal{L} = \bar{\psi}\ [\gamma \cdot D[A] + m(s)]\ \psi\,, \tag{14}$$

where the mass parameter depends on s, and vanishes at endpoints 0 and s_0: $m(0) = m(s_0) = 0$. The light quarks are "zero modes" that propogate in "our world" at $s = 0$ and s_0, one wall for each helicity. These new methods have shown early success [73,75] in preserving chiral symmetry, in terms of relations such as Eq. (13). In the lattice domain wall construction, the extra dimensions are just a convenience, but it cannot be denied that the technique bears an eerie resemblance to the brane constructions of modern string theory [76].

D. Hadrons and Instantons

Successes in the treatment of chiral symmetry in lattice QCD lead us naturally to the reemergence of instantons studies of hadron dynamics [77]. In the language of spontaneous symmetry breaking, the masses of hadrons are related to the generation of quark "condensates" $\langle \bar{q}q \rangle \sim \langle \bar{q}_R q_L \rangle + \langle \bar{q}_L q_R \rangle$ in the QCD vacuum, which couple (the perturbatively decoupled) left- and right-handed components of the quark field.

Instantons may be thought of as tunneling events between the inequivalent QCD vacuum configurations that are distinuished by different phases of the nonabelian gauge fields at infinity (even at zero energy). The instantons couple to the quarks, and the process of tunneling produces an effective $2N_f$-point interaction between their left- and right-handed components, of the general form

$$\mathcal{L} = G \, \frac{1}{8N^2} \left[\left(\bar{\psi} \tau^- \psi\right)^2 + \left(\bar{\psi} \tau^+ \gamma_5 \, \psi \,\right)^2 \right] \,, \tag{15}$$

even though the quarks are massless.

In the "instanton liquid" model, the nonperturbative soft gluon field is replaced by an ensemble of instantons and antiinstantons that couple left- and right-handed quark-antiquark pairs. Meson masses are produced as pairs "hop" between instantons, changing helicities as they go. In this model, mesons and baryons, with realistic spectra and matrix elements emerge, and provide a cross-check for lattice calculations [72]. at zero and at high temperatures.

IV. FACTORIZATIONS, EVOLUTIONS AND EFFECTIVE THEORIES

Every hard-scattering experiment includes a complete evolution all the way from short distance to long distance dynamics. Factorization allows us to organize the long distance dynamics, and thus to calculate perturbative short-distance dependence, and compare the results to experiment. The essence of factorization is to interpret long-distance information in terms of matrix elements in the underlying theory. For example, in the classic factorization in Eq. (6) for DIS structure functions of hadron h, $F^{(h)}(Q) = \sum_a C_a(Q/\mu) \otimes f_{a/h}(\mu)$, the quark ($a = q$) distributions may be interpreted as

$$f_{q/h}(x,\mu) = \int_{-\infty}^{\infty} \frac{dy}{4\pi} \, \mathrm{e}^{-ixP\cdot ny} \, \langle h(P)|\bar{q}(yn) n \cdot \gamma \, \Phi_n(y,0) q(0)|h(P)\rangle \,, \tag{16}$$

where n is a lightlike vector not in the direction of the hadron momentum P, and μ enters as the renormalization scale for the matrix element, which is ultraviolet divergent for $n^2 = 0$. The function $\Phi_n(y,0)$ is a path-ordered exponential of the gauge field,

$$\Phi_n(y,0) = P \exp\left[-ig \int_0^y d\lambda \; n \cdot A(\lambda n) \right] \,, \tag{17}$$

which makes the matrix element gauge invariant. The evolution of Eq. (3) may be thought of as a consequence of the renormalization properties of the nonlocal operators in $f_{a/h}$, which summarize an infinite set of twist-two matrix elements in the light-cone expansion. The same formalism is at the basis of heavy quark effective theory and of nonrelativistic QCD (NRQCD) [78], although generally with a finite sum over operators rather than a convolution. In a sense, the operator Φ_n in Eq. (16) plays the role of the heavy quark field in heavy quark effective theory, as a nonrecoiling source of gluons.

Over the past few years, factorization, analyzed in terms of effective operators, has been applied to multiscale hard scattering processes, with Sudakov [79,80] and Regge high-energy limits. The former refers to cross sections with large momentum transfer and low QCD radiation, the latter to low momentum transfer and essentially unlimited radiation. The Regge limit is related to the total cross section, while the Sudakov limit highlights its short-distance components.

One of the simplest examples of the Sudakov limit is the dijet cross section in $\mathrm{e}^+\mathrm{e}^-$ at fixed jet masses m_i^2, $i = 1, 2$. In the limit of light jets, $m_i^2 \ll Q^2$, the dijet cross section is related to the integrated thrust cross section by

$$\int_T^1 dT' \, \frac{d\sigma(T')}{dT'} = \int dm_1^2 dm_2^2 \, \frac{d\sigma}{dm_1^2 dm_2^2} \, \theta\left(1 - T - \frac{m_1^2 + m_2^2}{Q^2}\right) \,. \tag{18}$$

The cross section for "nearly-lightlike" jets in $\mathrm{e}^+\mathrm{e}^-$ satisfies a factorization [79]

$$\frac{d\sigma}{dm_1^2 dm_2^2} = \frac{C(Q/\mu, \beta_i)}{Q^6} \otimes \prod_{i=1}^{2} J_i(\mu, \beta_i) \otimes S(\mu, \beta_i)\,, \tag{19}$$

up to corrections suppressed by powers of m_i^2/Q^2, with β_i the 4-velocity of jet i. In (19), there is a double factorization, separating the dynamics of the jets, included in the functions J_i, from both the truely short-distance "coefficient" function C and from the dynamics of relatively low-energy partons emitted coherently by the jets and included in the function S. The "soft" function S is associated with a particularly interesting composite operator in QCD. S describes the emission of gluons whose wavelenths are so long that they cannot resolve the internal structure of the jets, and are thus generated by the product

$$W(0) = \Phi_{\beta_1}(\infty, 0)\, \Phi^\dagger_{\beta_2}(\infty, 0)\,, \tag{20}$$

where, in the notation of Eq. (17), the Φ's are ordered exponentials, and β_1 is the velocity of the quark jet, and β_2 of the antiquark jet.

We need not dwell on the nature of the convolutions denoted by $\otimes$ in Eq. (19), but the double factorization itself is adequate to imply a resummation [79] of double logarithms in $1-T$ [81],

$$\frac{1}{\sigma_{\rm tot}} \int_T^1 dT' \frac{d\sigma(T')}{dT'} = \exp\left[-2C_F \int_{1-T}^1 \frac{d\alpha}{\alpha} \int_{\alpha^2 Q^2}^{\alpha Q^2} \frac{dk_T^2}{k_T^2}\, \frac{\alpha_s(k_T^2)}{\pi} + \ldots \right]\,, \tag{21}$$

where corrections include fewer logarithms of $1-T$ in the exponent. These results are also related to the renormalization properties of the operators W in Eq. (20).

Another application of factorization and effective operators is to resummation in the "Regge" limit $s \to \infty$, t fixed, the "BFKL" regime for QCD. The BFKL equation may be derived from a "multiperipheral" reexpression of DIS factorization, Eq. (2) [1,82,83]:

$$\begin{aligned} F(x, Q^2) &= \int_x^1 \frac{d\xi}{\xi}\, C\left(\frac{x}{\xi}, \frac{\mu^2}{Q^2}\right) G(\xi, Q^2) + \mathcal{O}\left(1/Q^2\right) \\ &= \int d^2k_T\; c\left(\frac{x}{\xi'}, Q, k_T\right) \psi(\xi', k_T) \; + \mathcal{O}\left(1/\ln(1/x)\right)\,, \end{aligned} \tag{22}$$

where the k_T-dependent distribution ψ is related to the gluon PDF by [84]:

$$G(\xi, Q^2) = \int^Q d^2k_T\; \psi(\xi, k_T)\,. \tag{23}$$

In the second form of Eq. (22), the roles of longitudinal and transverse momenta have been reversed, and corrections are suppressed only by logarithmics of x, rather than powers of Q. In these terms, the BFKL equation,

$$\xi \frac{d\psi(\xi, k_T)}{d\xi} = \int d^2k'_T\; \mathcal{K}(k_T, k'_T) \psi(\xi, k'_T) = -\frac{\alpha_s N}{\pi^2} \int \frac{d^2k_T}{(k_T - k'_T)^2} \left[\tilde\psi(\xi, k_T) - \frac{{k'_T}^2}{2k_T^2} \tilde\psi(\xi, k'_T) \right]\,, \tag{24}$$

with $\tilde\psi \equiv (1/k_T^2)\psi$, describes the evolution of $\psi(\xi, k_T)$ in ξ. The same equation may also be derived from the renormalization of ordered exponentials, like Eq. (17). Indeed, distributions of exponentials integrated over a transverse distribution [82]

$$\int d^2x_T\; \rho(x_T)\; P\; \exp\left[-ig \int dx^+ A^-(x^-, x_T) \right] \tag{25}$$

are being studied to develop "unified" effective theories to describe the variety of evolution equations, including DGLAP, BFKL and others [83].

V. BFKL, DIFFRACTION AND COLOR DYNAMICS

A. BFKL at NLO

What's so special about the BFKL equation? It addresses the total cross section in gauge theory in terms of its fundamental quanta, the quarks and gluons. Now it is not entirely obvious that such a project will work, ultimately, in an asymptotically free theory with confinement, but if it does, it will say something fundamental about field theory. In an older language, this would be a theory of the "pomeron" [85]. Also, in the language of the parton model, BFKL appears to predict that, as we evolve to low $x \leftrightarrow$ high $s_{\gamma^* N}$ at fixed Q, we reach a region of high parton density at nearly fixed (actually slowly diffusing [86]) virtuality. This is a new "intermediate" regime of QCD, between perturbative and hadronic phases [87]. It is relevant to the screened "plasma" state, which we hope to encounter at RHIC.

The solutions to the lowest order BFKL equation are of the form

$$\tilde{\psi} \sim x^{-\omega} \left(\frac{k_T^2}{\mu^2}\right)^{-i\nu-1/2} . \tag{26}$$

The largest permissible value of ω gives the dominant low-x behavior, which is found to be

$$\nu = 0, \quad \omega = \omega_0 \equiv 4N \ln 2\, (\alpha_s/\pi) . \tag{27}$$

From the kinematic relation in DIS, $s_{\gamma^* N} = Q^2(1-x)/x$, the low-$x$ behavior of $\psi(x, k_T)$ determines the large-s behavior of the $\gamma^* N$ total cross section, for which the lowest-order BFKL result (27) gives

$$\psi \sim \text{``}\sigma_{\rm tot}\text{''} \sim s^{4N \ln 2\, (\alpha_s/\pi)} . \tag{28}$$

This is a derivation of Regge-like behavior for the total cross section from perturbative QCD, and, because it involves the exchange of no overall quantum numbers, may be considered as a perturbative model for the pomeron.

Where should one look for BFKL behavior in experiment? Suggestions include correlations in dijet and rapidity-gap cross sections in DIS and $p\bar{p}$, and in non-DGLAP evolution in DIS at low x and moderate Q^2. In the first case, there may be hints in the dijet data from HERA and in the comparison of jet correlations at 630 and 1800 GeV [11]. In the later case, the strong rise in x of the structure functions $F_i(x, Q^2)$ cannot be sustained indefinitely, since, by (28), this would eventually violate unitarity bounds for the cross section. Before this happens, interference between partons, or "shadowing", which is absent in both DGLAP and BFKL evolutions, must begin to set in [88]. One of the much-discussed data presentations of the past year, Fig. 6 from the ZEUS collaboration, shows the transition between perturbative and nonperturbative behavior in a particularly suggestive form [89].

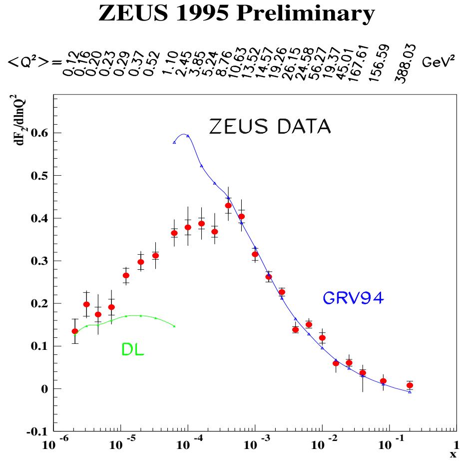

FIG. 6. Slopes of F_2 as measured by ZEUS.

1998 was the year of the NLO BFKL kernel, the year in which the decade-long project of computing the next-to-leading order kernel in Eq. (24) bore fruit [90]. At NLO, $\mathcal{K}$ is fairly complicated, but the effect on ω in Eq. (27) is simple enough,

$$\omega^{(NLO)} = \omega_0 \left[1 - 6.6 \left(\frac{N\alpha_s}{\pi} \right) \right] . \tag{29}$$

Now 6 and a half is not by itself a large number, but the size of this corretion nevertheless presents a challenge, because unless α_s is quite small, the second term may overwhelm the first and lead to an unrealistic falling cross section. In addition, it has been observed that the NLO kernel even implies a non-Regge behavior at high orders [91]. Interesting responses to these challenges were discussed in [92]. Evidently, the NLO fruit of BFKL will be an acquired taste, but the coming year surely promises intensive work and further clarification

B. Diffraction and Diffractive PDFs

Through the optical theorem, the total cross section is closely related to elastic and diffractive scattering amplitudes [14]. In diffractive DIS, we can relax inclusivity and thus probe QCD dynamics in the final state, while retaining a large momentum transfer. Convenient variables to describe diffraction are

$$x_{\mathcal{P}} = \frac{M_X^2 + Q^2}{W^2 + Q^2} \qquad \beta = \frac{x}{x_{\mathcal{P}}} , \tag{30}$$

where M_X is the mass of an observed system X moving in the "current" (*i.e.* photon) direction, and $W^2 = s_{\gamma^* s}$. $x_{\mathcal{P}}$ is the longitudinal momentum transfer from the nucleon to X, and β is the equivalent fractional momentum of a parton in the (hypothetical) exchanged "pomeron".

Diffractive events are typically defined by a large gap in rapidity between X and the elastically-scattered, or diffractively-excited, proton (or its low-mass fragments), which has experienced a momentum transfer t. In these terms, a fully differential cross section is

$$\frac{d^4\sigma}{Q^2 d\beta dx_{\mathcal{P}} dt} = \frac{2\pi\alpha_{\rm EM}^2}{\beta Q^4} \, \left(1 + (1-y)^2\right) \, F_2^{D(4)} . \tag{31}$$

In the simplest diffractive processes X consists of a single vector boson. In this case, the relevant amplitude is factorized in terms of off-diagonal PDF's, related for small x to the gluon distribution, which behaves as $G(x) \sim x^{-\lambda}$. The resulting cross section is proportional to $G^2(x)$, and hence increases with W, as

$$\frac{d\sigma^D}{dM_X} \sim \frac{1}{M_X^4} \, W^{4\lambda(M_X)} . \tag{32}$$

There is compelling evidence for this behavior, with a value of λ increasing with M_X, suggesting once again a Regge-like behavior reminiscent of BFKL. This connection has yet to be completely explored.

High-Q^2 DIS diffractive cross sections may be factored [93] using specifically diffractive PDFs, also referred to as fracture functions [94],

$$F^D = \sum_a \, C_a^D \otimes f_a^D . \tag{33}$$

Phenomenological fits to $f_{q,g}^D$'s have been carried out [95], and lead to predictions whenever a factorization like Eq. (33) applies [96]. It is important to realize, however, that because diffractive PDFs are not fully inclusive, they depend on the details of evolution into the final state, in particular the likelyhood of the target proton staying together. This probability cannot be expected to be the same in $\mathrm{p\bar{p}}$ cross sections, where the fragments of two initial-state hadrons pass through each other, as in DIS, where only a single hadron is involved. And indeed, studies [97] have shown that diffraction at the Tevatron is much less likely than would be suggested by a direct generalization of Eq. (33) to this case with universal diffractive PDFs.

Nevertheless, double-diffraction (double rapidity gap) jet production is seen at Tevatron, and the jets show a standard parton-parton E_T-dependence [97], indicating that the short-distance process is independent of the long-distance evolution. This suggests that another factorization is possible in this case, and may shed light on diffractive dynamics.

C. Color Dyamics

Diffractive processes are naturally interpreted in terms of color-singlet exchange in the t-channel. Large momentum transfer processes, however, are described perturbatively in terms of single-gluon exchange, carrying octet quantum numbers. Over the past few years, there has been progress in understanding the relation between these two pictures.

Although color is not observable, representations $(1, q = 3, g = 8 \ldots)$ are, at least in principle. In NRQCD, a factorization that includes the mixing of operators with differing color content has already led to valuable insights and phenomenological successes [78]. For high-energy processes, rapidity gaps were long ago suggested by Bjorken as an ideal arena to study color exchange, with singlet exchange expected to produce an excess of events with very low interjet multiplicity.

It is clear, however, that it is not possible to separate short- from long-distance color exchange uniquely, since gluons of all momenta carry the same color content, and models in which the color content of the final state is determined at the longest distances have had success [98]. At the same time, energy flow into regions between two high-p_T jets is senstive to all time scales between $1/|p_T|$ and $1/\Lambda_{\rm QCD}$. Energy flow at the shorter time scales is both perturbative and sensitive to the color content in the t-channel. This observation led [99] to an analysis of dijet cross sections in terms of energy flow Q_c into the interjet region, in the range $\Lambda_{\rm QCD} \ll Q_c \ll \sqrt{-t}$. The cross section at fixed Q_c may be factorized, and logarithms of Q_c^2/t resummed. This behavior is found from the renormalization properties of composite operators that are products of ordered exponentials, which generalize Eq. (20) to the $2 \to 2$ scattering of partons with color exchange (labelled here by f),

$$W^{\rm f}_{\beta_j \alpha_j}(0) = \prod_{i=3}^{4} \Phi_{n_i}(\infty, 0)_{\beta_i,\kappa_i} T^{\rm f}_{\kappa_4 \ldots \kappa_1} \prod_{i=1}^{2} \Phi_{n_i}(0, -\infty)_{\kappa_i,\alpha_i} \tag{34}$$

where $T^{\rm f}$ is a matrix that couples the color of the incoming and outgoing ordered exponentials Φ_{n_i}, representing the active partons of the hard scattering. These operators mix under renormalization and induce an evolution that tracks mixing in the color space as the scattering particles (including $q\bar{q}$, qg, gg) evolve from short to long distances. This analysis offers a new set of predictions for p_T, energy and rapidity dependence of gap events, which can be tested at Run II of the Tevatron, and at the LHC.

VI. POWER CORRECTIONS

There has been considerable interest in power corrections to infrared safe quantities. As noted in Sec. 2 above, such power corrections are quite important in the phenomenology of jet cross sections and event shapes in e^+e^- annihilation [22,32]. Behind this work is a hypothesis, that it is not necessary to model the details of hadronization to parameterize leading corrections to perturbation theory, and a hope, that plausible parameterizations inspired by perturbation theory will lead to useful insights at the perturbative-nonperturbative interface [31]. The hypothesis seems to be correct; whether the hope will be realized remains to be seen, but there are preliminary indications that it might be [100].

How does perturbation theory imply nonperturbation corrections? In the calculation of any IR safe quantity at NLO, we always encounter integrals of the general form

$$I(\alpha,p) = f(x)\ \alpha_s(Q^2)\ \int_0^Q dk\ k^p\,, \tag{35}$$

with $p > -1$, where k may be thought of as a gluon momentum scale, and $f(x)$ represents the remaining (IR finite) dependence. Now in many cases (see Eq. (21) for example) we can argue (or derive) that higher order corrections modify (35) to

$$I^{(\mathrm{resum})}(\alpha,p) = f(x) \int_0^Q dk\ \alpha_s(k^2)\ k^p \ldots . \tag{36}$$

This result shows that the perturbative reexpansion in terms of $\alpha_s(Q)$ is asymptotic, with high orders that grow as $n!$ at large n. This information is encoded in the singularity of the perturbative expression for $\alpha_s(k^2)$ at $k = \Lambda_{\mathrm{QCD}}$. We can further reinterpret this behavior by means of a Borel transform, but the inverse transform will not be unique in any case. Taking a more practical approach, we retain the perturbative factor $f(x)$ and the perturbative range in the k integral in Eq. (36), and simply replace the lower end of the integral, $k < \mu_1$, for some fixed μ_1, by a parameter $\alpha_p(\mu_1)$,

$$\int_0^{\mu_1} dk\ \alpha_s(k^2)\ k^p \rightarrow \mu_1^{p+1}\ \alpha_p(\mu_1)\,, \tag{37}$$

Since the overall integral in Eq. (35) behaves as Q^{p+1} this is automatically a power correction. Clearly the value of this approach depends on the assumption that $\alpha_p(\mu_1)$ is in some sense universal [100]. Evidently, this is true approximately, and this method has found applications in models for DIS higher twist [101]. As the notation suggests, the parameter α_p' is naturally thought of as a reflection of a universal, nonperturbative low-scale running coupling. This is clear in (36) above, where it is a moment of the lowest-order running $\alpha_s(k^2)$. It has been argued that the relation is more general, and that higher orders of $\alpha(k^2)$ may be incorporated into a reconstructed effective coupling, defined through dispersion relations [102,103].

In interpreting these developments, it is important to keep in mind that the values of higher-twist parameters cannot be defined independently of perturbation theory [100], and that they will change as new orders, or resummations, are computed. A striking example of this effect was illustrated by the NNLO analysis of [17], which reduced the size of higher twist contributions, relative to those found in fits based on NLO.

Another interesting application of these ideas is to resummed event shapes, as in Eq. (21), where a replacement like Eq. (37) leads to a simple shift [104] in the thrust ($T = 1 - t$) distribution, which vanishes as $1/Q$,

$$\frac{d\sigma_{\mathrm{PT}}(t)}{dt} \rightarrow \frac{d\sigma_{\mathrm{PT}}(t-\lambda/Q)}{dt} + \mathcal{O}\left(\frac{1}{(tQ)^2}\right)\,, \tag{38}$$

with λ a constant related to α_0 in (37). In a somewhat more general approach, we may once again factorize soft gluon emission into the region between the two jets, and derive a convolution expression for the cross section [105],

$$\frac{d\sigma_{\mathrm{PT}}(t)}{dt} \rightarrow \int_0^{tQ} d\epsilon\ f(\epsilon)\ \frac{d\sigma_{\mathrm{PT}}(t-\epsilon/Q)}{dt} + \mathcal{O}\left(\frac{1}{tQ^2}\right)\,, \tag{39}$$

where $f(\epsilon)$, a "shape function", has a field-theoretic interpretation [105,106], which involves the composite operator of Eq. (20). $f(\epsilon)$ is Q-independent and summarizes all $1/(tQ)^n$ corrections, while Eq. (39) reduces to (38) with the replacement $f(\epsilon) = \delta(\epsilon - \lambda)$. A fit [105] to the extensive thrust data based on this formalism, and the perturbative resummation of [81], is shown in Fig. 7. Given the discussion of the foregoing section, by following this line of reasoning we may hope to relate event shapes in e^+e^- annihilation to energy flow in hadronic hard-scattering cross sections. This relation remains unexplored, although it is certainly related to multiplicity and correlation studies of the final states in jet events [107].

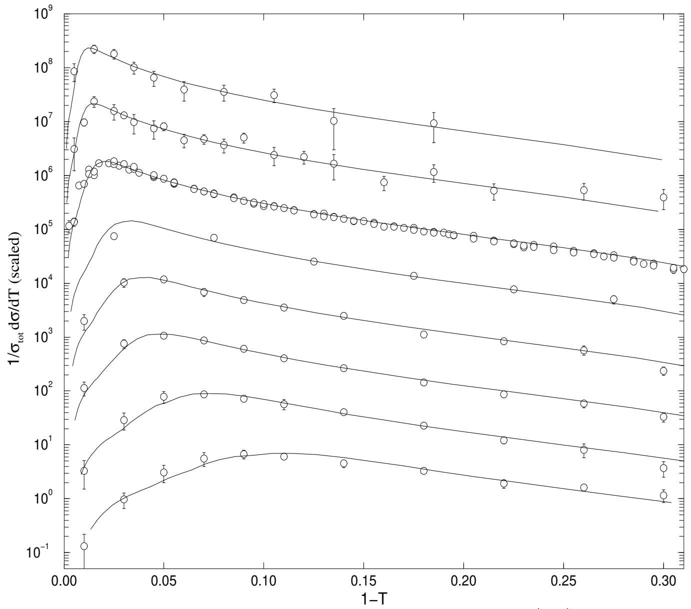

FIG. 7. The comparison of thrust data with a fit based on Eq. (39), at energies (from bottom to top): $Q/\mathrm{GeV} = 14, 22, 35, 44, 55, 91, 133, 161$.

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VII. QCD AT HIGH TEMPERATURE AND BARYON NUMBER

I have already mentioned the path from perturbative QCD to high parton density through BFKL evolution. Certainly, these considerations are made more interesting by the pending turn-on of the RHIC accelerator at Brookhaven, at which nuclei will be collided at unprecented energies. This development has led to a fresh look at QCD in "extreme" conditions, long of relevance to studies of the early universe.

Color Superconductivity. At high enough density and temperature, the long-distance interactions that lead to confinement in the normal, hadronic phase of QCD are screened, and in some sense nonsinglet degrees of freedom are freed. In fact, QCD is expected to have a possibly quite rich phase structure in the plane of temperature (T) and baryon density (B).

An exciting exploration of these features of the theory is the fresh look at the long-standing conjecture of color superconductivity at large B and low T, in the light of the instanton liquid model referred to above in Sec. IIID above. It was realized in Ref. [108] that the four-fermion effective Lagrange density in Eq. (15) produces an attractive potential between quarks, which can lead to a condensate of quark Cooper pairs at the Fermi surface, in a manner analogous to superconductors of electric current. In fact, the condensate can be driven by gluon exchange, but the energy gaps produced by instantons are much larger in most, but not all [109], of parameter space. Although more likely to be relevant for neutron stars than nuclear collisions, these interesting considerations were clearly inspired by RHIC physics, which has led to an efflorescence of studies of the QCD B/T plane [110,111].

Energy Loss. As a final example, I will very briefly refer to some recent considerations on a topic of direct interest to RHIC and hadron-nucleus scattering, energy loss of fast partons in dense media. High-energy partons travelling through matter (partons or hadrons) will scatter and radiate, and their evolution into the final state will be modified in some way. In sufficiently inclusive processes, these effects are (perhaps surprisingly) higher twist [112]. It is of interest, however, to also look at changes in radiation at transverse momentum scales set by the medium, rather than a hard scattering. Along these, lines, recent work on the energy loss [113] experienced by a quark travelling through a medium over length (L) has identified a QCD analog of the famous Landau-Pomeranchuk effect in QED. This work analyzes "induced" gluon radiation at $\Lambda \ll k_T \ll \langle Q \rangle$, where $\langle Q \rangle$ is the typical momentum transfer in a projectile-medium scattering. If the Debye length of the medium is short enough in a hot, dense medium, Q could be perturbative.

The amplitude for an emission $q \to q + g$ at impact parameter b is denoted $f(b,t)$. The effect of the scatterings is a diffusion in b,

$$f(b,t) \sim b \left\langle Q^2 \frac{d\sigma}{dQ^2} \right\rangle \exp \left[-i \text{ const } \left\langle Q^2 \frac{d\sigma}{dQ^2} \right\rangle b^2 t + \ldots \right] , \tag{40}$$

with known corrections to the Gaussion. For large $\langle Q \rangle$, the cross section involves the interference between amplitudes where the radiation occurs at different positions along the path through the medium, leading to a power spectrum in path length and frequency given by

$$\omega \frac{dI}{d\omega\, dz} \sim \alpha_s \int_0^L dt \left(1 - \frac{t}{L} \right) \int d^2 b\; f(b,t)\; f(b,0) \,. \tag{41}$$

This results in an energy loss per unit length that grows with length (!):

$$-\frac{dE}{dz} \sim \frac{\alpha_s N_c L}{4\lambda} \left\langle Q^2 \frac{d\sigma}{dQ^2} \right\rangle , \tag{42}$$

at mean free path λ, which is a diagnostic for $\left\langle Q^2 \frac{d\sigma}{dQ^2} \right\rangle$, potentially able to distinguish the composition of the medium, whether hadrons, Debye-screened plasma or something else. Here again, the transition from energy loss at relatively low transverse momentum, to more inclusive cross sections, should be an interesting one [114].

VIII. CONCLUSION

The central conclusion of this little review is the variety and vitality of the work itself. Beyond this, the nature of our knowledge of QCD is such that important ideas and techniques only require the possibility that they can be tested to receive further theoretical development. The ongoing experiments at LEP, HERA and the Tevatron have already transformed perturbative QCD into a truely quantitative discipline. The pending RHIC accelerator has inspired creative theoretical developments. A strong QCD component to future high-energy projects is sure to be richly rewarded by insights into quantum field theory .

ACKNOWLEDGEMENTS

I would like to thank the organizers of DPF 99 for the invitation to participate in a vibrant meeting. I am indebted to Werner Vogelsang for invaluable help. This work was supported in part by the National Science Foundation, under grant PHY9722101.

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