STATUS OF PERTURBATIVE QCD

John Ellis
CERN, Geneva, Switzerland

Introduction

Theorists now believe they know how to make asymptotic predictions for (almost) all hadronic processes at large momentum transfers. To an increasingly total degree, perturbative QCD[1] can now be applied to any 'hard' scattering process previously the domain of the parton model[2]. The basic physical picture of the parton model is retained, but perturbative QCD enables us to compute novel modifications, subasymptotic corrections, and interrelations between different processes. Much of the theoretical work underlying this explosive expansion of applications of QCD has taken place in the two years since the last lepton-photon symposium[3]. This accounts for much of the difference in character between this and theoretical talks at previous symposia, although many prophesies can be found there[4,5], particularly in the 1975 talk of Polyakov[4].

Another reason for a major change in the character of these talks has been the increasing pace of attempts[6,7] to make quantitative confrontations of the 'classical' QCD predictions and experimental results[8-10], particularly in the domain of deep inelastic scattering. The first -- and largest -- part of this talk discusses and assesses these different 'classical' tests of QCD, such as σ_{tot} ($e^+e^- \to$ hadrons)[11], the charmonium model[12], and deep inelastic scattering structure functions[13]. The second part of this talk discusses the 'asymptotically free parton model'[14,15] and its applications to inclusive final-state hadrons in deep inelastic collisions[16], the Drell-Yan process[17], jets[18,19], and the interactions of real photons at large momentum transfers[20]. The third part of this talk introduces the application[21] of perturbative QCD to exclusive 'hard' processes such as form factors and wide-angle scattering, which has excited much interest in the last few months. Section 4 mentions some new directions which go beyond the ritual re-summation of leading and non-leading logarithms in perturbative QCD. Section 5 summarizes prospects and problems for experimental tests and theoretical applications of perturbative QCD.

1. Classical Applications

Our ability to make high-energy predictions from field theories of the strong interactions rests on the renormalization group[22]. The quantitative predictivity of perturbative QCD rests on its unique[23] property of asymptotic freedom[24]: the renormalization group shows us that its effective coupling at a momentum scale Q decreases, $\alpha_s(Q^2) \to 0$ as $Q^2 \to \infty$. Traditionally, the tool for applying the renormalization group and asymptotic freedom to physical hadronic processes has been the operator product expansion[25]. Most 'classical' applications of perturbative QCD, such as $\sigma_{tot}(e^+e^- \to$ $\to$ hadrons) and deep inelastic scattering, depend on these three ingredients.

Solving the renormalization group equations for QCD in leading order, we find[24] the asymptotic freedom property

$$\alpha_s(Q^2) \approx \frac{12\pi}{(33-2f)\ln Q^2/\Lambda^2} \tag{1}$$

where f is the number of 'operational' quark flavours ($4m_q^2 << |Q^2|$) and Λ is an integration constant which represents the 'scale of the strong interactions'. The renormalization group and operator product expansion enable us to expand physical quantities in powers of $\alpha_s(Q^2)$:

$$[\text{Phys.}] = [\text{Phys.}]_0 \left[1 + O(\alpha_s) + O(\alpha_s^2) + \cdots\right] \tag{2}$$

To go beyond the leading-order prediction $[\text{Phys.}]_0$ in Eq. (2) generally requires going beyond the leading-order formula (1) for $\alpha_s(Q^2)$. In the next order[26]

$$\alpha_s(Q^2) \approx \frac{12\pi}{(33-2f)\ln Q^2/\Lambda^2}\left[1 - \frac{\beta_1 \ln\ln Q^2/\Lambda^2}{\beta_0^2 \ln Q^2/\Lambda^2}\right] \tag{3}$$

where $\beta_1 = 102 - 38/3\ f$, $\beta_0 = 11 - 2/3\ f$. There is a new integration constant $\tilde{\Lambda}$ which is again an hadronic 'scale', but comparison of formulae (1) and (3) shows that $\Lambda \neq \tilde{\Lambda}$. Indeed, we see that the non-leading-order term $O(1/\ln Q^2)^2$ has a dependence involving $\ln\ln Q^2$, which is inconsistent with the simple Λ parametrization of Eq. (1). Unless care is taken, the parametrization of subasymptotic phenomena in QCD using a Λ parameter is meaningless, and there is no reason why the Λ's extracted from the subasymptotic corrections to different processes should agree. However, if one works consistently using higher-order formulae such as (3), one can in principle interrelate the Λ parameters measured in different processes.

A further problem is that beyond leading-order, different prescriptions for renormalizing QCD give different results, unless one is able to sum over all orders of perturbation theory, which is likely to be never possible. We are therefore stuck with ambiguities corresponding to the arbitrariness in choice of renormalization prescription at any given order of perturbation theory.

Much of the art in comparing higher-order calculations with experiment therefore lies in the choice of a renormalization prescription which makes for rapid convergence in calculations of the process (or finite set of processes) being considered. Even better, one may look for predictions which are independent of the choice of renormalization prescription. In general, it should be remembered that different prescriptions (as well as processes) are characterized by different Λ parameters, which may, however, be interrelated in the context of higher-order calculations as we will see quantitatively later on. Let us first see these remarks illustrated by more explicit examples.

1.1 $\sigma_{total}(e^+e^- \to \gamma^* \to \text{hadrons})$

The naîve parton model prediction from Fig. 1a for the physical quantity $R \equiv \sigma(e^+e^- \to \gamma^* \to \text{hadrons})/\sigma(e^+e^- \to \gamma^* \to \mu^+\mu^-)$ is[2]

$$R = 3\sum_{q=1}^{f} e_q^2 \tag{4a}$$

The leading QCD correction of Fig. 1b -- ignoring complications[24] due to the transition from space-like

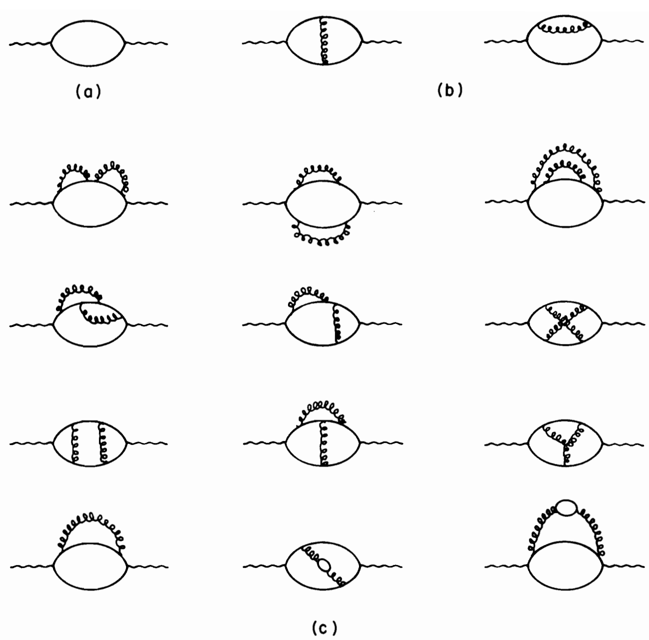


Fig. 1 QCD diagrams for e^+e^- annihilation: a) in zeroth order; b) in second order; c) in fourth order in the strong gauge coupling.

$Q^2 < 0$ (where the operator product expansion and renormalization group analysis strictly applies) to time-like $Q^2 > 0$, which should be all right away from new quark thresholds -- gives[11]

$$R = 3 \sum_{q=1}^{f} e_q^2 \left[1 + \frac{\alpha_s(Q^2)}{\pi} \right] \tag{4b}$$

The next-to-leading QCD correction (Fig. 1c) to (4a) has recently been computed[28], and gives

$$R = 3 \sum_{q=1}^{f} e_q^2 \left[1 + \frac{\alpha_s(Q^2)}{\pi} + \begin{Bmatrix} 7.35 - 0.442 f \\ 1.98 - 0.115 f \end{Bmatrix} \left(\frac{\alpha_s(Q^2)}{\pi} \right)^2 \right] \tag{4c}$$

The second-order correction depends on the renormalization prescription. The top line in (4c) refers to the 'minimal subtraction' scheme (MS), where we only remove the poles in $1/\varepsilon$ encountered in the dimensional regularization procedure[29]. The bottom line in (4c) refers to the 'minimal subtraction' scheme $\overline{MS}$, where all associated factors of $(\ln 4\pi - \gamma_E)$ have also been removed[30]. The corresponding Λ parameters [cf. Eq. (3)] are related by

$$\Lambda_{\overline{MS}} = \Lambda_{MS} \exp(4\pi - \gamma_E)/2 \approx 2.66\, \Lambda_{MS} \tag{5}$$

Comparing theory and experiment, we find[31] the results given in Table 1 for a centre-of-mass energy Q = 6 GeV, and $\Lambda_{\overline{MS}}$ = 0.5 GeV motivated by the deep inelastic experiments to be discussed later.

We see that with a good choice of renormalization prescription -- in this case the $\overline{MS}$ scheme[30] -- the perturbation expansion converges well. The present experimental error is much larger than the theoretical uncertainty -- it is of the same magnitude as the first-order QCD correction, while the second-order QCD correction is an order of magnitude smaller. The principal experimental uncertainty is systematic, and the precision of present theoretical calculations makes worth while an experimental effort to reduce this error to or below the present statistical error. An over-all error of 3% seems to be an attainable and desirable target.

1.2 Quarkonium Decays

Another 'classical' (in the sense that it was generally accepted at the time of the previous symposium) application of perturbative QCD is to the decays of heavy quarkonia[12]. We visualize these decays as taking place through a heavy $Q\bar{Q}$ annihilation at a very small distance $O(1/m_Q)$ into a collection of light quanta: gluons g, light quarks q, and photons. The simplest instance is the paraquarkonium ${}^1S_0(Q\bar{Q}) \to 2g$ (Fig. 2a), which is identified in leading order with the total hadronic rate. The leading-order QCD prediction for this rate relative to ${}^1S_0(Q\bar{Q}) \to 2\gamma$ is[12]

$$\frac{\Gamma({}^1S_0 \to \text{hadrons})}{\Gamma({}^1S_0 \to \gamma\gamma)} = \frac{\Gamma({}^1S_0 \to gg)}{\Gamma({}^1S_0 \to \gamma\gamma)} = \frac{2}{9 e_Q^2} \left(\frac{\alpha_s(4m_Q^2)}{\alpha} \right)^2 \tag{6}$$

We have used the asymptotically free coupling constant at the 'scale' of the bound state, although it is not clear until (or even after) a non-leading calculation whether this is the appropriate coupling to use. A

Table 1

Theory and experiment for $e^+e^- \to$ hadrons

Source of contribution	naive parton model	leading QCD corrections	QED vacuum polarization	Quark mass corrections	next-to-leading QCD corrections	Total
Value of ΔR	3.33	0.32	0.13	0.088	-0.029	3.84
			Experimental value[31]			4.17 ±0.09 (stat.) ±0.42 (syst.)

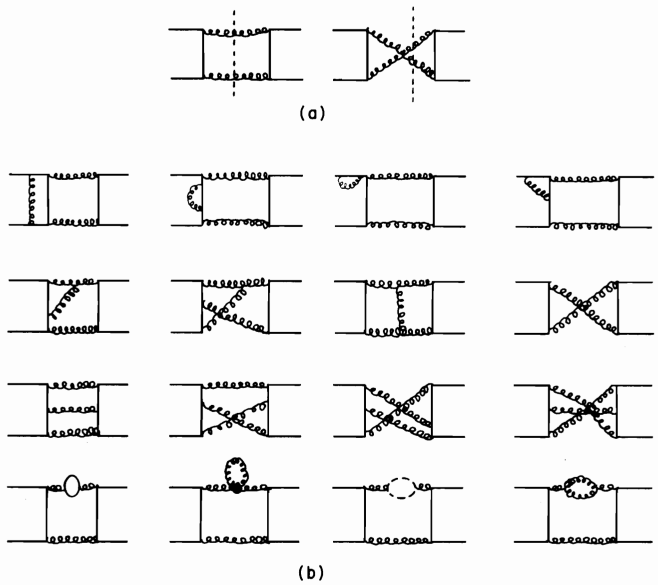


Fig. 2 QCD diagrams for paraquarkonium decay: a) in lowest order; b) in next-to-lowest order.

calculation of the first-order QCD radiative corrections (Fig. 2b) to Eq. (6) has now been done[32], with the dramatic result

$$\frac{\Gamma({}^1S_0 \to \text{hadrons})}{\Gamma({}^1S_0 \to \gamma\gamma)} = \frac{2}{9e_Q^2}\left(\frac{\alpha_s(4m_Q^2)}{\alpha}\right)^2\left[1 + \begin{Bmatrix} 22.14 \\ 16 \end{Bmatrix} \frac{\alpha_s(4m_Q^2)}{\pi}\right] \tag{7}$$

where the top (bottom) line refers to the MS ($\overline{\text{MS}}$) renormalization scheme, respectively. Included in the left-hand side of Eq. (7) are gg, ggg, and $gq\bar{q}$ final states. In this case we see that the QCD radiative corrections (7) are enormous whichever of the two (MS or $\overline{\text{MS}}$) schemes we use.

What lessons do we draw from this débâcle? It certainly seems that the conventional lowest-order calculations[33] for other onium decays (${}^3S_1 \to ggg$ or $gg\gamma$, ${}^3P_{0,2} \to gg$, ${}^3P_1 \to gq\bar{q}$) are invalid for charmonium, very questionable for bottomonium, and perhaps only barely applicable to toponium decays. It is clear that more calculations of radiative corrections to onium decays are in order. Corrections to the ${}^3S_1 \to ggg$ process are probably an order of magnitude more complicated than those for ${}^1S_0 \to gg$. Corrections to ${}^3S_1 \to gg\gamma$ should be intermediate in complication, and topical because of present measurements[34] of the decay rate and photon spectrum for $J/\psi \to \gamma + X$. Corrections to ${}^3P_{0,2} \to gg$ might be relatively simple to calculate, and also very topical because of the apparent[35] discrepancy between the experimentally deduced[36] and lowest-order theoretical ratios of their hadronic decays:

$$\frac{\Gamma({}^3P_0 \to \text{hadrons})}{\Gamma({}^3P_2 \to \text{hadrons})} \neq \frac{15}{4} \tag{8}$$

Another interesting problem is whether, despite the large magnitude of the radiative corrections, the expected two- and three-jet structures of onium final states[37] (${}^1S_0$, ${}^3P_0$, ${}^3P_2 \to 2$ jets, ${}^3S_1 \to 3$ jets) may nevertheless survive. After some more of these onium radiative corrections have been calculated, it may even turn out that there is a better choice than $\alpha_s(4m_Q^2)$ to use in the onium decay rate formulae, which may consistently give more rapidly convergent perturbation series than using the MS or $\overline{\text{MS}}$ schemes for onium decay calculations.

1.3 Deep Inelastic Scattering

This is the truly 'classic' application of perturbative QCD underwritten[38] by all the paraphernalia of the operator product expansion[25] and the renormalization group[22]. Moments of the deep inelastic structure functions

$$M_N(Q^2) = N\int_0^1 dx\, x^{N-1} \begin{Bmatrix} W_1(x,Q^2) \\ x^{-1} W_2(x,Q^2) \\ W_3(x,Q^2) \end{Bmatrix} \tag{9}$$

are related to the matrix elements of local operators of definite spin, whose asymptotic behaviour is explicitly calculable[13] using the renormalization group and asymptotic freedom. (The prefix N indicates that the moments should be Nachtpersonnalized[39] to continue the definite spin projection down to finite Q^2. The predictions apply to moments of the full deep inelastic 'cross-section', so all elastic and quasi-elastic final states should be included[40].) For non-singlet combinations of structure functions (e.g. F_2^{ep-en}, $F_3^{\nu N + \bar{\nu} N}$) the predictions for the moments (9) are[13]

$$M_N(Q^2) \approx A_N\left(\ln Q^2/\Lambda^2\right)^{-d_N} : d_N = \frac{4}{33-2f}\left[1 - \frac{2}{N(N+1)} + 4\sum_{j=2}^{N}\frac{1}{j}\right] \tag{10}$$

while for singlet combinations there are two leading terms:

$$M_N(Q^2) \approx A_N^+\left(\ln Q^2/\Lambda^2\right)^{-d_{N^+}} + A_N^-\left(\ln Q^2/\Lambda^2\right)^{-d_{N^-}} \tag{11}$$

The predictions (10) and (11) for the moments of the structure functions can be inverted[41] to give the evolution with Q^2 of the structure functions directly. We will return to this later. For the moment we consider proposed tests of the moment behaviours (10) and (11), which are the direct predictions of the traditional approach to perturbative QCD.

Two simple tests of the behaviour (10) expected for the non-singlet moments in QCD have been proposed[8]. On the basis of formula (11) one expects that

$$\ln M_N(Q^2) = (\text{constant}) + \frac{d_N}{d_M}\ln M_M(Q^2) \tag{12}$$

and that

$$\left[M_N(Q^2)\right]^{-1/d_N} \propto \ln Q^2/\Lambda^2 \tag{13}$$

These two formulae test different aspects of the theory. In Eq. (12) the slope parameters d_N/d_M just reflect the spin of the gluon coupled to the quark[42], assuming that the scaling violation comes predominantly from the lowest-order bremsstrahlung of a gluon ('gluestrahlung'). This can however only be strictly justified in an asymptotically free theory, of which QCD is the only example. In lowest order, all vector gluon theories yield

$$\frac{d_N}{d_M} = \left[\frac{1-\frac{2}{N(N+1)}+4\sum_{j=2}^{N}\frac{1}{j}}{1-\frac{2}{M(M+1)}+4\sum_{j=2}^{M}\frac{1}{j}}\right] \tag{14}$$

while all scalar gluon theories give[43]

$$\frac{d_N}{d_M} = \left[\frac{1-\frac{2}{N(N+1)}}{1-\frac{2}{M(M+1)}}\right] \tag{15}$$

The second prediction (13) tests the asymptotic freedom $\alpha_S \sim 1/\ln Q^2$ of QCD: in a fixed-point theory where the coupling constant $\alpha_S \to$ constant $\alpha^* \neq 0$, the right-hand side would grow like a power of Q^2.

Many questions have been raised about the reliability and significance of tests of these moment prepredictions of QCD. What about higher orders of QCD perturbation theory? One expects them to modify the prediction (10) by a factor

$$\left[1+O\left(\frac{\alpha_s}{\pi}\right)+O\left(\frac{\alpha_s}{\pi}\right)^2+\dots\right] \tag{16}$$

The complete $O(\alpha_S/\pi)$ corrections have been calculated for both singlet and non-singlet moments[44], and one can discuss their effects on the naive lowest-order predictions (10) and (11) at presently accessible Q^2. What about higher-twist effects[45]? These are expected to modify the naive leading-order perturbative QCD prediction by a factor

$$\left[1+O\left(m^2/Q^2\right)+O\left(m^2/Q^2\right)^2+\dots\right] \tag{17}$$

where m is some typical hadronic mass scale. These corrections are difficult to compute, and the magnitude of their effects at present Q^2 is largely unknown[46]. Faced by experimental problems as well as these theoretical questions about the forms of the moments, many people suggest[45] that one should perform direct analyses of the Q^2 evolution of the structure functions themselves. We will return later to this type of analysis: for the moment we concentrate on the primary QCD predictions [formulae (10) to (13)].

1.4 Higher Orders of Perturbation Theory

The first point we will study is the effect of higher orders of perturbation theory[44]. In keeping with the general principles mentioned earlier on, we will be looking for quantities with a good (rapidly convergent) perturbation expansion, and quantities which are independent of the renormalization prescription used. How do the proposed tests (12) and (13) meet these criteria?

It has been shown[47,48] that the second-order perturbation theory terms in (16) enable us to compute the corrections to the lowest-order values d_N/d_M of these slopes. We find[47]

$$\ln M_N(Q^2) = (\text{constant}) + \ln M_M(Q^2)\frac{d_N}{d_M}\left[1+\left(\frac{\tilde{d}_N}{d_M}-\frac{\tilde{d}_M}{d_N}\right)\frac{\alpha_s}{4\pi}+\dots\right] \tag{18}$$

where the $\tilde{d}_N$, $\tilde{d}_M$ are two-loop anomalous dimensions. The coefficients C_{MN} of α_S in Eq. (18) are given in Table 2 for some experimentally interesting values of M,N.

Table 2

Higher-order corrections to ln moment slopes

M	N	C_{MN}
2	4	0.42
4	6	0.21
6	8	0.15

We see that the coefficients C_{MN} of α_S are small, indicating a good perturbation expansion for $\ln M_N$ versus $\ln M_M$ plots[8-10], with corrections to the slopes which are O(10)% at moderate values of Q^2. The lowest-order (QCD1) and higher-order (QCD2) slopes are shown in Fig. 3 together with data from ν and eN scattering experiments[49]. The data are clearly consistent with the QCD predictions, but the errors are too large to discriminate between the lowest- and higher-order QCD predictions.

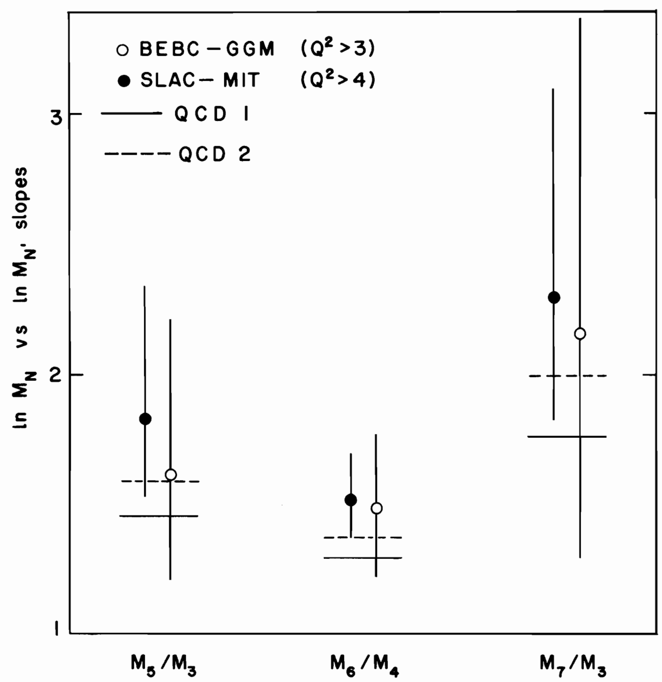


Fig. 3 Leading (QCD1) and next-to-leading order (QCD2) predictions for logarithmic moment slopes, compared with data taken from Barnett[49].

As for the plots of $[M_N(Q^2)]^{-1/d_N}$ [Eq. (13)], it has been shown[47] that they should be linear in $\ln Q^2$ to a good approximation, but that the extraction of Λ^2 is very uncertain. In fact the effective Λ parameter introduced in Eq. (10) is inadequate because it does not take account of the expected $\ln\ln Q^2/\ln Q^2$ [cf. Eq. (3)] corrections to the leading-order moment predictions, and in addition the subasymptotic corrections are N-dependent, so that the effective Λ parameter should vary with N[50]. The simplest renormalization prescription-independent way of doing this in a manner consistent with the expected $\ln\ln Q^2/\ln Q^2$ corrections is[51] to parametrize the moments by

$$\left[M_N(Q^2)\right]^{1/d_N} \approx \frac{\overline{M}_N}{\left(\ln Q^2/\Lambda_N^2\right)}\left[1-\frac{\beta_1}{\beta_0^2}\frac{\ln\ln Q^2/\Lambda_N^2}{\ln Q^2/\Lambda_N^2}\right] \tag{19}$$

Not only are the Λ parameters dependent on N, they also depend[51] on which structure function is being analysed, e.g. F_2 or F_3. All the different Λ_N may be related to the basic scale parameter of α_s, which we take to be $\Lambda_{\overline{MS}}$. Shown in Table 3 below are values of the Λ_N for different N and structure functions, all expressed in units of $\Lambda_{\overline{MS}}$.

Table 3

Theoretical values of Λ_N

N	2	5	10
F_2	1.34	1.95	2.54
F_3	0.98	1.83	2.48

Figure 4 shows some values of Λ_N extracted from experiment[52], and curves illustrating the expected N dependence of Λ [53]. The Λ_N come together as $N \to \infty$ but differ at finite N, so that one should *not* be overjoyed if neutrino and electron/muon experiments get identical values of Λ_N! The "BEBC/GGM" and "CDHS" values of Λ_N are much closer together than the published values. This is partly because the expected $\ln \ln Q^2/\ln Q^2$ behaviour in the moment parametrization (19) was not included originally, and partly because of differences in the assumptions under which the data were originally analysed, which are given in Table 4 [52,54]. The uniform analysis assumptions used here are indicated by asterisks.

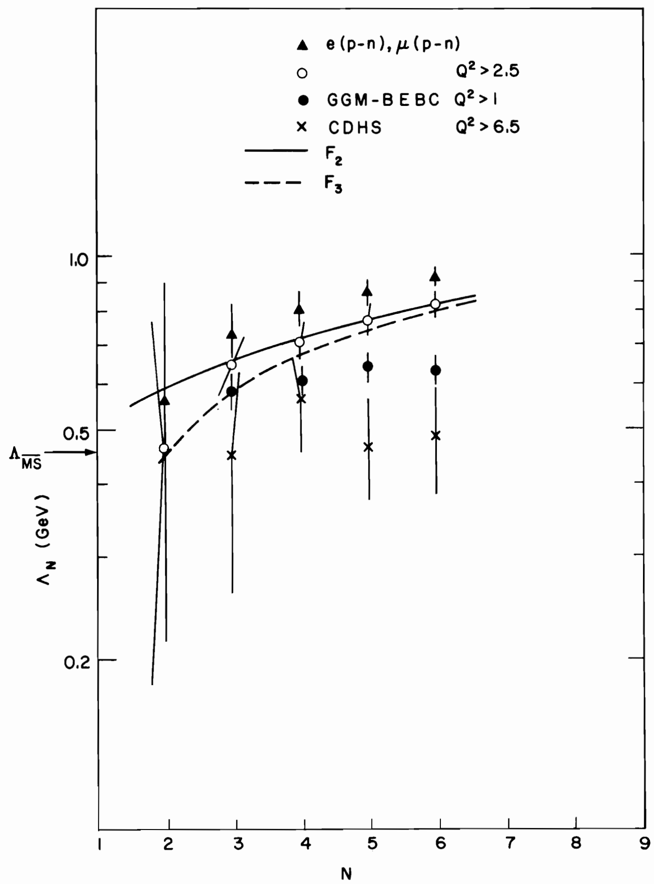


Fig. 4 Values of Λ_N defined by Eq. (19) extracted from experiment[52] and compared with the N-dependences expected from QCD[51]. The vertical scale on the theoretical curves is arbitrary: they may be moved up and down at will. Shown on the vertical axis is the corresponding value of $\Lambda_{\overline{MS}}$ [30].

Standardizing the analysis on each of these four points has the effect of reducing the previous apparent discrepancy, as does the inclusion of $\ln \ln Q^2/\ln Q^2$ correction terms. As for the electron-muon data in Fig. 4, the crossed lines indicate an analysis without the $\ln \ln Q^2/\ln Q^2$ term; the circles indicate our estimate of the likely effect of including it[52]. Since the Λ_N are plotted on a logarithmic scale, and since theory does not predict the absolute scale but only the ratios, the theoretical curves may be moved together up or down. The electron-muon data nicely reproduce the trend expected theoretically, while the neutrino data are somewhat discrepant but perhaps not grossly so. (Notice, however, that the CDHS Λ_6 is about 40% different from the electron-muon Λ_6, whereas they are expected to differ by < 5%; all is not rosy.) The value of $\Lambda_{\overline{MS}}$ corresponding to the plotted theoretical Λ_N curves is shown on the vertical axis. The different experiments probably correspond to the range

$$0.35\ \mathrm{GeV} < \Lambda_{\overline{MS}} < 0.5\ \mathrm{GeV} \tag{20}$$

So far we have discussed exclusively non-singlet structure functions which should have the simple asymptotic form (10) (where A_N is related to moments V_N of the valence quark distributions) rather than the more complicated form (11) (where the $A_N^{\pm}$ are related to moments G_N and S_N of the gluon and singlet quark distributions). The more complicated structure (11) reflects the fact that there are two singlet sets of operators, made up of gluons and quarks[13], corresponding to two different singlet parton distributions. These intercommunicate by the $q\bar{q}$ and GG pair creation diagrams of Fig. 5b, as well as the basic gluestrahlung diagram of Fig. 5a. The $d_N^{\pm}$ are generally positive, except that $d_2^- = 0$. This means that $M_2(Q^2)$ -- which in parton language measures the longitudinal momentum fraction carried by quarks -- is a fixed constant asymptotically. This asymptotic constancy reflects an equilibrium[42] between the diagram generating gluons from quarks (Fig. 5a)

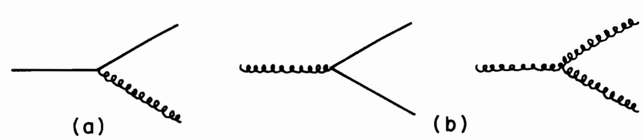


Fig. 5 Vertices controlling scaling violations in lowest-order QCD.

Table 4

Analysis differences between BEBC/GGM and CDHS

	Number of flavours	Fermi motion	Radiative corrections	Quasi-elastics
BEBC/GGM	3	Corrected*	Not included	Included exactly*
CDHS	4*	Uncorrected	Included*	Parametrized approximately

*Assumptions used in computing the Λ_N in Fig. 4.

and that generating quarks from gluons (Fig. 5b). The value of the equilibrium point depends on the field theory being considered, and is different for QCD and an Abelian vector gluon theory[55]. Shown in Fig. 6 are present values of $M_2(Q^2)$ for νN and μp scattering, which are both consistent with falling towards the asymptotic QCD value rather than rising towards the asymptotic Abelian value.

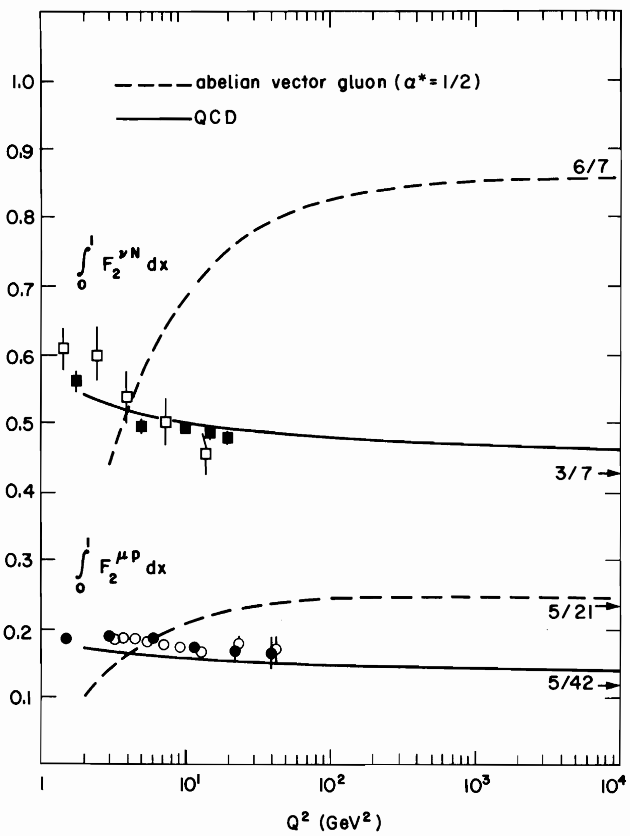


Fig. 6 Moments of $F_2(X,Q^2)$ compared[55] with predictions of QCD and an Abelian vector gluon theory.

When we come to higher singlet moments, there is no asymptotically constant solution, and tests of the QCD prediction (11) are more subtle. An approach pioneered by the BEBC-GGM group[8], followed by Duke and Roberts[53] and put into its most recent form by Perkins[56], goes as follows. Introduce two quantities

$$T_N \equiv \frac{S_N}{V_N} \quad , \quad R_N \equiv \frac{G_N}{S_N} \tag{21a}$$

and denote

$$\chi_N^{0,\pm} \equiv T_N^{\,d_N,\, d_{N\pm}} \tag{21b}$$

then if we plot $\chi_N^0(T_N/T_{N^0}) - \chi_N^+$ vertically versus $\chi_N^- - \chi_N^+$ horizontally (where $T_{N^0} \equiv T_N$ at some reference momentum Q_0^2), we should get a straight line

$$f(d_N, d_{N\pm}) - R_N\, g(d_N, d_{N\pm}) \tag{21c}$$

where f and g are known functions and R_N was defined in Eq. (21a). A plot[56] of this type for $M_3(Q^2)$ using

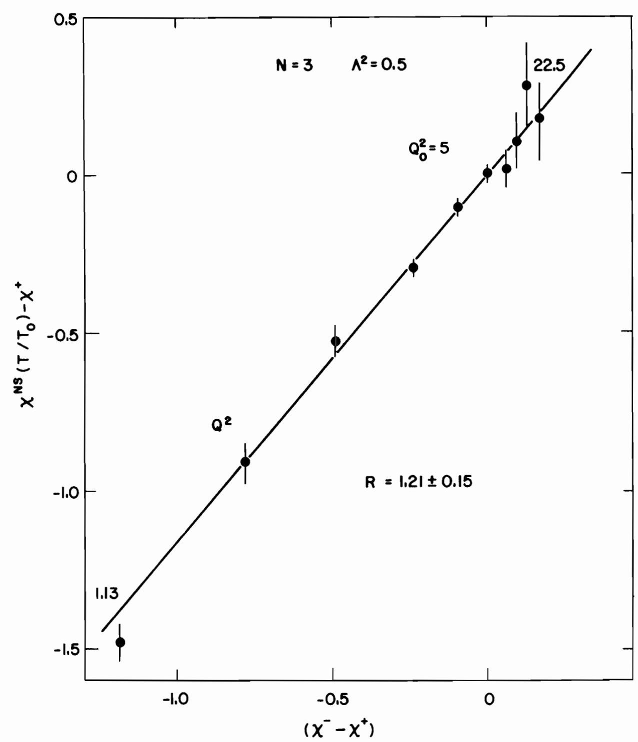


Fig. 7 Singlet moments of $F_2^{e(p-n)}$ plotted[56] as described [Eq. (21)] in the text, and a QCD straight line fit.

electron and muon data for $1.13 \leq Q^2 \leq 22.5$ GeV^2 is shown in Fig. 7. The data are consistent with the expected linear behaviour and with

$$R_3(Q_0^2 = 5\,\mathrm{GeV}^2) = \frac{G_3}{S_3} = 1.21 \pm 0.15 \tag{22a}$$

while a similar analysis (not shown) of $M_5(Q^2)$ yields

$$R_5(Q_0^2 = 5\,\mathrm{GeV}^2) = \frac{G_5}{S_5} = 1.52 \pm 0.45 \tag{22b}$$

These values of the gluon moments are not absurd[57], and indicate that the singlet structure function data, as well as the non-singlet data discussed earlier, are not glaringly inconsistent with QCD perturbation theory.

1.5 Higher-twist Effects

We now turn to the thorny question whether the perturbative analysis makes any sense, or is invalidated by higher-twist effects. (Note that 'higher-twist' is theoretical jargon[13] for $O(m^2/Q^2)$ corrections to leading-order QCD predictions -- to my knowledge it has no rational explanation.) It is a fact[45,58] that all deep inelastic scaling violations so far observed can be fitted by higher-twist effects alone without aid from the QCD logarithms (10) and (11). This ambiguity in the interpretation of scaling violations applies in equivalent forms to direct analyses of scaling violations in deep inelastic structure functions and to moment analyses. Shown in Fig. 8 are perturbative QCD and higher-twist fits[45] to BEBC-GGM data between Q^2 of 1 and 70 GeV^2. They are equally good[59].

Theoretically we expect[46] the higher-twist effects on moments to be relatively more important at large N (structure functions near X = 1):

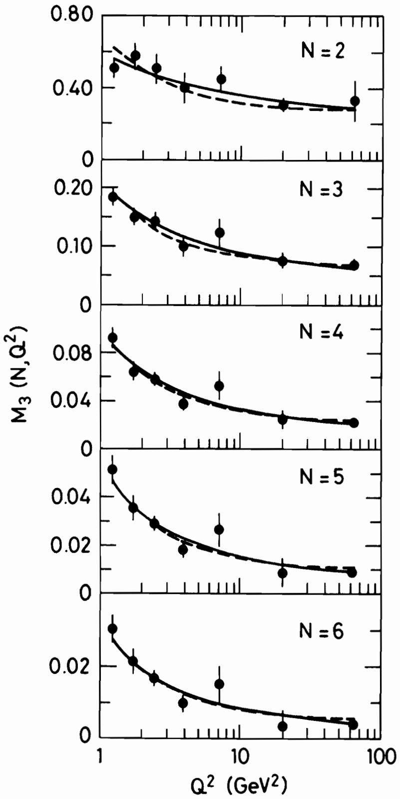


Fig. 8 Fits to moments of $xF_3^{\nu N}(X,Q^2)$ using perturbative QCD (solid lines) and higher-twist effects (dashed lines) taken from Ref. 45.

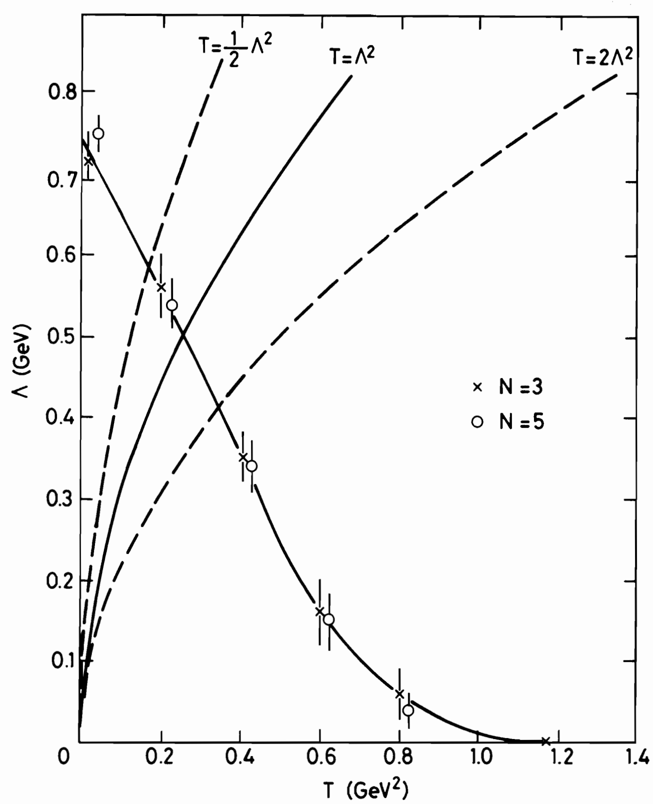


Fig. 9 Fits to the third and fifth moments of $xF_3(X,Q^2)$ (BEBC-GGM data) using[56] two parameters Λ and T.

$$\frac{\Delta M_N}{M_N} \underset{N\to\infty}{\sim} N\frac{T}{Q^2} \tag{23}$$

where T is a dimensional higher-twist parameter that cannot be calculated at the present level of theoretical understanding. From the form of (23) we see that neglect of higher-twist effects is only justifiable[40,46,60] for

$$N \ll Q^2/T \tag{24}$$

so that QCD perturbation theory should be a good approximation only *non-uniformly* in N and Q^2. In order to see which N obey the criterion (24) at any given value of Q^2, we need to know the value of T. Theorists[46] have suggested that

$$T = O(\langle p_T\rangle^2) = O(0.1 \text{ to } 0.2)\,\mathrm{GeV}^2 \tag{25}$$

where $\langle p_T\rangle$ is a typical hadronic transverse momentum, presumably $O(1/R_N)$, where R_N is the nucleon radius. A good fit to the totality of electron, muon, and neutrino scaling violations requires[45,56]

$$T \simeq 1.3\,\mathrm{GeV}^2 \quad : \quad \frac{\Delta M_N}{M_N} \simeq \frac{1.3\,(N-3/2)}{Q^2} \tag{26}$$

If one tries a combined fit using T and the leading-order Λ as free parameters, as shown in Fig. 9, one finds[56] that their values are highly correlated with any ratio of T/Λ^2 between 0 and ∞ giving an acceptable fit. However, since both T and Λ^2 measure closely related hadronic scales, it somehow seems unreasonable that either $T \gg \Lambda^2$ or that $\Lambda^2 \gg T$. If we apply the restriction $T = O(\Lambda^2)$, as indicated by the curves for $T = (\frac{1}{2},1,2)\Lambda^2$ in Fig. 9, we find that the preferred value of Λ is reduced by a factor of the order of 1.5. It remains to be seen whether this is a reasonable guess at the magnitude of higher-twist effects. We theorists should strive very hard to get a better theoretical understanding of the magnitude of higher-twist effects, and we hope that more precise deep inelastic data will help to resolve the present Λ versus T ambiguity. At the moment this is the greatest uncertainty in the analysis of deep inelastic scaling violations.

1.6 Direct Analysis of Structure Functions

Is the moment analysis we have been discussing up to now the best way of testing perturbative QCD in deep inelastic structure functions? The theoretical advantage of moments is that there are precise numerical predictions for their behaviour: slope of ln M_5 versus ln M_3 = 1.456 + 0.27 α_s + ... , but experimentally they are difficult to measure (sensitivity to elastics and quasi-elastics, badly measured high values of X, etc.) It has been suggested[45,61,62] that we should return to the old method[41] of directly analysing the scaling violations in structure functions. We now have more refined tools than we had previously for doing this, in particular the Altarelli-Parisi[63] equations which express the QCD effect on the quark and gluon distributions in a direct, mathematically elegant way. For a valence quark distribution,

$$Q^2\frac{\partial}{\partial Q^2} q^V(x,Q^2) = \frac{\alpha_s(Q^2)}{2\pi}\int_x^1 \frac{dy}{y} P_{q\to q}\left(\frac{x}{y}\right) q^V(y,Q^2) + \cdots \tag{27}$$

where $P_{q\to q}(z)$ is a known function

$$P_{q\to q}(z) = \frac{4}{3}\left[\frac{1+z^2}{(1-z)_+} + \frac{3}{2}\,\delta(z-1)\right] \tag{28}$$

Analagous formulae exist for gluon and singlet quark distributions [with analogous functions $P_{q\to g}(z)$, $P_{g\to q}(z)$, and $P_{g\to g}(z)$] and the $O(\alpha_s^2)$ corrections in (27) are explicitly known. The formula (27) and its friends are logically *equivalent* to the moment predictions (10) and (11), with

$$\int_0^1 dz\, P_{q\to q}(z) \;\propto\; d_N \tag{29}$$

Because of this equivalence, there are *equivalent* ambiguities and uncertainties in the treatment of higher-order and higher-twist effects, etc. However, Eq. (27) enables us to compute directly the scaling violations at some (X_0,Q_0^2) in terms of the structure function at $X \geq X_0$ and some $Q^2 < Q_0^2$. This means that we do not have to know the structure function at low $X \leq X_0$, and in fact the forms of $P_{q\to q}(z)$ and of the valence quark distributions for $X > 1/4$ are such that the principal contribution to $Q^2(\partial/\partial Q^2)^2 q^V(X_0,Q^2)$ comes from X only slightly $> X_0$. This makes the Altarelli-Parisi[63] equation (27) a convenient one to work with. We can either use it in the differential form (27)[61] or in an integrated form[62]:

$$q^V(X_0,Q_0^2) = \int_{X_0}^1 dX\, K(X,X_0;Q^2,Q_0^2)\, q^V(X,Q^2) \tag{30}$$

where the kernel K is known explicitly to $O(\alpha_s^2)$.

Several comparisons of formulae (27) or (29) with experimental data have been made[61,62]. As an example, see Fig. 10 taken from a paper[61] which uses the original differential form (27). The shaded region indicates the trend of scaling violation in $Q^2(\partial/\partial Q^2)^2 F_2^{ep-en}(X,Q^2)$ ($Q^2 = 3.2$ GeV^2) to be expected from lowest-order QCD with a leading-order $\Lambda \sim 0.45$ to 0.55 GeV. The crosses are extracted from SLAC data, and the qualitative agreement is remarkable. This method of analysis is certainly promising: it remains to be seen whether it can be made as interestingly quantitative as the moment analyses.

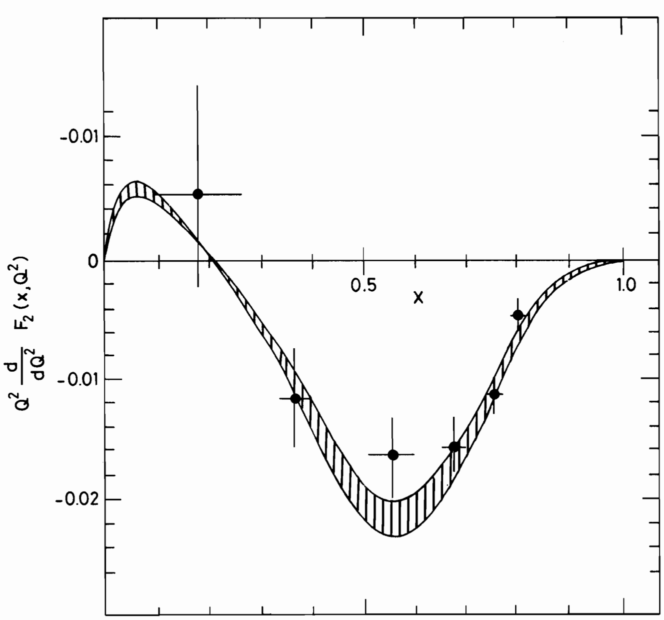


Fig. 10 SLAC-MIT data at $Q^2 = 3.2$ GeV^2 are used[61] to deduce (shaded region) the logarithmic scaling violation to be expected[63] at that Q^2 as a function of X, and compared with the experimental scaling violations.

After all these analyses are done, what is the preferred value of Λ? As was emphasized above, stating a precise value has meaning only when higher orders are computed and included in the analysis. Then we find that the effective Λ parameters are different for different processes. For example:

$$\Lambda_2^{F_2} = 1.34\,\Lambda_{\overline{MS}}\;;\;\Lambda_2^{F_2} = 1.37\,\Lambda_2^{F_3}\;;\;\Lambda_{\overline{MS}} = 2.66\,\Lambda_{MS} \tag{31}$$

Table 5 lists the methods and results of several different second-order analyses of deep inelastic structure functions. Some are moment analyses; some are direct structure function analyses. They quote values for a number of differently defined Λ parameters, which have all been translated into equivalent values of $\Lambda_{\overline{MS}}$ using (31) and other analogous formulae.

Table 5

Λ parameters extracted from deep inelastic structure functions

Type of analysis	Type of Λ estimates	Corresponding value of $\Lambda_{\overline{MS}}$
$\mu p-\mu n$, F_2^{ep-en} moments[52]	$\Lambda_n^{F_2}$	0.45 ± 0.1?
$XF_3^{\nu N}$ moments[52]	$\Lambda_{\overline{MS}}$	0.5 ± ?
$XF_3^{\nu N}$ moments[52]	$\Lambda_n^{F_3}$	0.45 ± 0.1? (BEBC-GGM) 0.35 ± 0.2? (CDHS)
$F_2^{eN,\mu N}$ directly[62]	Λ_{MS}	0.7 ± 0.3
$F_2^{e(p-n)}$ directly[62]	Λ_{MS}	0.53 ± 0.16
$XF_3^{\nu N}$ directly[62]	Λ_{MS}	0.46 ± 0.21

It is perhaps surprising that there should be so much consistency between the different analyses, particularly bearing in mind the gross differences in systematics between direct and moment analyses. A reasonable mean of all these values is

$$\Lambda_{\overline{MS}} \sim 0.5\,GeV \tag{32}$$

with an unknown error. The analyses quoted above ignore higher-twist effects. In line with the discussion above, the effect of a 'reasonable(?)' magnitude of higher-twist effects may be to reduce $\Lambda_{\overline{MS}}$ by up to a third. Some people might ask what is the point of determining $\Lambda_{\overline{MS}}$ very precisely, since the effects of varying it are apparently almost imperceptible. But remember that the mass of the proton is probably roughly proportional to Λ, and some day we hope to calculate it! A more immediate hope is to calculate in grand unified theories the proton lifetime τ which is roughly proportional to Λ^4 for smallish variations in Λ. The most recent calculations[64] suggest that $\Lambda_{\overline{MS}} = 0.5$ GeV [Eq. (32)] corresponds to $\tau \sim 2 \times 10^{30}$ years, with perhaps an error of ±1 in the exponent. Another ±1 in the exponent comes from uncertainty in $\Lambda_{\overline{MS}}$; and people about to commit several years to an underground life would clearly like to see this uncertainty reduced.

2. The Asymptotically Free Parton Model

The 'classical' applications of asymptotic freedom and perturbative QCD were those sanctified by the renormalization group[22] and the operator product expansion[25]. These processes were also in the domain of the naive parton model[2], which however was also applied to many processes where the renormalization group/operator product expansion approach was apparently not applicable[65]. In the last two years it has been learnt[14] in QCD perturbation theory how to justify, modify, and extend the parton model in the bulk of these applications. Let us first state the general result for an inclusive hard-scattering process, and then briefly review the elements of its derivation.

Let us consider, as an archetypal hard-scattering process, hadron-hadron scattering to produce particles at large p_T as illustrated in Fig. 11. The naive parton model[2] would describe this process using distributions $P_i(X_i)$ of incoming partons, Born scattering cross-sections, and final-state fragmentation functions $D_j(Z_j)$:

$$\sigma = P_1(X_1)\,P_2(X_2)\,\sigma^{\mathrm{Born}}_{12\to34}\,D_3(Z_3)\,D_4(Z_4) \tag{33a}$$

(In any physical process there will in general also be convolutions over initial and/or final distributions, and a sum over the collisions of different parton types.) Perturbative QCD takes into account radiative corrections to the basic hard-scattering process, coming for example from gluestrahlung (cf. Fig. 5a) and pair creation (cf. Fig. 5b). These QCD radiative corrections have several effects. One is to introduce a scale (Q^2) dependence into the initial parton distributions and final-state fragmentation functions. Another is to introduce the effective running coupling (1) into the formulae for the Born cross-sections. Another is to add to the leading-order cross-section calculable higher-order terms corresponding to a power series in $\alpha_s(Q^2)$. Thus Eq. (33a) is replaced by

$$\begin{aligned}\sigma = {} & P_1(X_1,Q^2)\,P_2(X_2,Q^2)\\ & \times \sigma^{\mathrm{Born}}_{12\to34}(\alpha_s(Q^2)) \qquad + O(\alpha_s(Q^2))\\ & \times D_3(Z_3,Q^2)\,D_4(Z_4,Q^2)\end{aligned} \tag{33b}$$

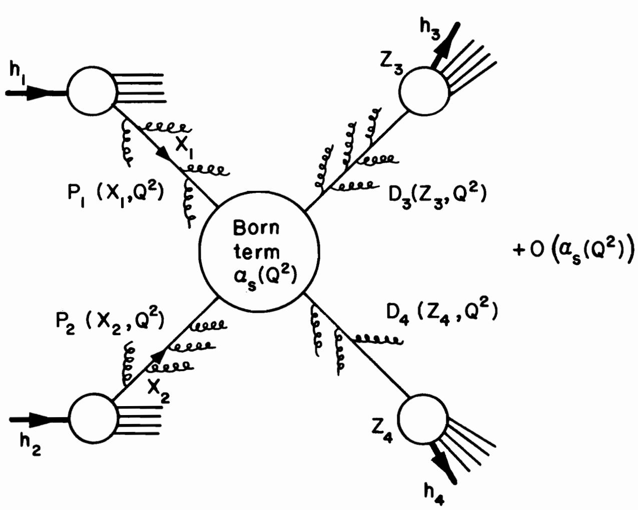


Fig. 11 The generic asymptotically free parton model[14] structure for large p_T hadron-hadron collisions.

The Q^2 scale is generally of the order of the large kinematic invariants which are (ideally) taken to ∞ in a constant ratio in the hard-scattering process (e.g. $Q^2 = sX_1X_2$ in hadron-hadron collisions[66]. The initial parton distributions obey the Altarelli-Parisi[63] equations (27), while the fragmentation functions obey very similar (transposed) equations. These Q^2 dependences come predominantly from collinear gluestrahlung or pair creation, as indicated in Fig. 11. The effective coupling $\alpha_s(Q^2)$ in the Born cross-section emanates from vertex corrections, while hard gluestrahlung and wide-angle pair creation are parts of the $O[\alpha_s(Q^2)]$ correction terms in Eq. (33b). The initial (or final) hadrons in Eq. (33b) may be replaced by real photons[20], while pairs of hadronic legs may be replaced by virtual intermediate bosons (γ^*, W^* or Z^*) or lepton pairs, and analogous results apply.

Many groups of authors have participated in the justification of the asymptotic free parton model (33b)[14]. A key observation, first made in the context of simple low-order diagrams[67], was that the large logarithms $\ln Q^2/p_i^2$ (where p_i is an initial or final parton momentum) factorized, external leg by external leg, and being universal could be absorbed into the definitions of the Q^2-dependent distributions $P_i(X_i,Q^2)$ and $D_j(Z_j,Q^2)$. These could then be taken from deep inelastic scattering or $e^+e^- \to h + X$, respectively, and used in a multitude of other applications. Particularly elegant are diagrammatic analyses[68] of the leading and non-leading logarithms of QCD which reveal how the renormalization group results [Eqs. (10) and (11)] are built up diagram by diagram, and in particular they have the jet structure of the hadronic final states[18,19]. An interesting way of restating the results uses cut vertices[15], whose definition and renormalization strongly recall those of local operators, and suggests a basis for the asymptotically free parton model which is as rigorous as the operator product expansion in perturbation theory. Of course, in all these analyses it is assumed that non-perturbative (e.g. confinement) effects do not alter the conclusions of the perturbative analysis. This assumption has been vigorously criticized[65], but an equivalent assumption has always been made in the renormalization group/operator product expansion analysis of deep inelastic scattering. However, it remains true that it has not yet been proved[69] that non-perturbative effects do not mess up the beautiful formula (33b). Anyway, let us now summarize a few of the interesting applications.

$e^+e^- \to h + X$

The inclusive hadron cross-section in e^+e^- annihilation can be used[70] as the definition of the fragmentation function $D(Z,Q^2)$ (cf. Fig. 12), which violates scaling in a calculable way analogous to deep inelastic structure functions. These violations are small at high (PETRA/PEP) energies, and may best be seen by comparing low-energy (SPEAR-DORIS) and high-energy data, or by

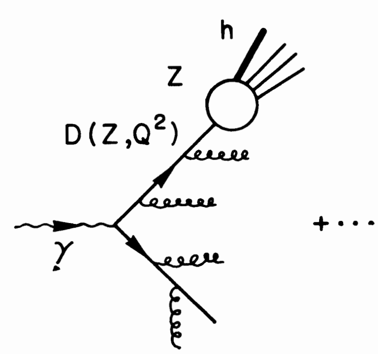


Fig. 12 QCD prediction for $e^+e^- \to h + X$.

looking more systematically at inclusive hadron production in the centre-of-mass region between 1.5 and 3.5 GeV. It would be interesting to see Z-moments of these inclusive hadron cross-sections[71].

$\ell + h \to \ell' + h' + X$

The cross-section for this process (Fig. 13) can be written as

$$\frac{d^2\sigma}{dz\,dx} = q(x,Q^2)D(z,Q^2) + O(\alpha_s(Q^2)) \qquad (34)$$

We expect violations of scaling, and the $O[\alpha_s(Q^2)]$ terms violate factorization[16] and generate hadrons at large p_T [72] relative to the incoming virtual intermediate boson (γ^*, W^*, or Z^*).

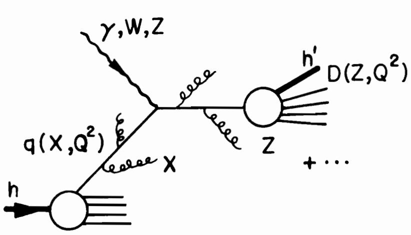


Fig. 13 QCD prediction for $\ell + h \to \ell' + h' + X$.

$h_1 + h_2 \to \ell^+\ell^- + X$

The simple-minded Drell-Yan cross-section[17] of Fig. 14 is modified to become

$$\sigma = q_1(x_1,Q^2)\bar{q}_2(x_2,Q^2)\left(\frac{4\pi\alpha^2}{3Q^2}\right)\frac{1}{3} + O(\alpha_s(Q^2)) \qquad (35)$$

In this case the violations of scaling are small in the kinematic region $4 < M_{\ell^+\ell^-} \equiv Q < 8$ GeV where most of the data exist, but the subasymptotic $\alpha_s(Q^2)$ corrections cause a *large renormalization*[73] of the cross-section by a *factor* O(2).

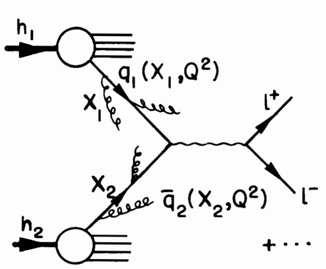


Fig. 14 QCD prediction for $h_1 + h_2 \to (\ell^+\ell^-) + X$.

$h_1 + h_2 \to h + X$

The large p_T cross-section of Fig. 11 is given already in Eq. (33b). In this case a satisfactory description of the present-day large p_T data may need all the scaling violations $[P_i(X_i,Q^2),\ D_j(Z_j,Q^2),\ \alpha_s(Q^2)]$ available there, as well as substantial non-perturbative initial and final (fragmentation) state $\langle p_T \rangle$[74]. It seems reasonable to expect a large renormalization of this cross-section analogous to that for the Drell-Yan process (35), but the relevant computations have not yet been done.

Jets

One can easily compute, in the framework of the asymptotically free parton model, the cross-sections for producing multiple jets in hard-scattering processes, e.g. $e^+e^- \to 3$ jets[18] (Fig. 15), $\ell + h \to \ell' + 2$ forward

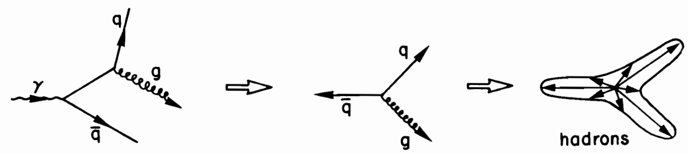


Fig. 15 QCD prediction[18] for three-jet events in e^+e^- annihilation due to $e^+e^- \to q\bar{q}g$.

jets + 1 backward jet[72], $h_1 + h_2 \to 3$ large p_T jets, etc.[75]. All one does in Eq. (33b) is discard the final-state fragmentation functions and put in a higher-order Born term, e.g. for a $2 \to 3$ reaction such as $q + q \to q + q + g$ in large p_T processes. The fundamental extra-jet process is $e^+e^- \to q\bar{q}g \to 3$ jets[18] due to wide-angle gluestrahlung.

Photons

'Hard' processes with an initial or final real photon can be considered in a very analogous way to that of the treatment of partons and hadrons[20]. There are contributions where the photon has a direct interaction (Fig. 16a), so that the initial parton distribution (or

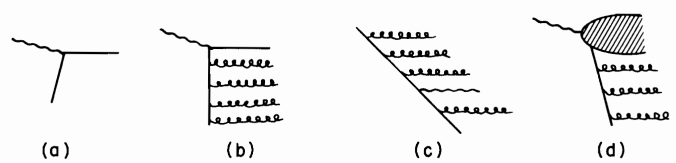


Fig. 16 a) The point-like photon; b) its hard structure function; c) its hard fragmentation function, and d) the soft part of its structure function, as seen[20] in perturbative QCD.

fragmentation function) in Eq. (33b) is removed, and the Born cross-section involves an incoming (or outgoing) real photon leg, e.g. for deep inelastic Compton scattering[76] $\gamma + q \to \gamma + q$, and its QCD analogue[77]. There are also hard processes probing the point-like distributions of quarks and gluons in a photon[78] (Fig. 16b) and analogously point-like fragmentation functions of quarks and gluons into photons (Fig. 16c), which are absolutely calculable and grow $\propto \ln Q^2$. Finally there are soft 'hadronic' contributions (Fig. 16d) to these distributions which fall with the conventional hadronic powers of $\ln Q^2$. This analysis gives a complete description of photons in 'hard' processes which amounts to an understanding of the old puzzle about when, and why, the photon acts in a point-like way.

After this brief résumé of some typical hard processes, we will now turn to a more detailed discussion of some of them which are of most immediate phenomenological interest.

2.1 Leptoproduction of Hadrons

As mentioned above, quark and gluon fragmentation functions should obey analogues[79] of Altarelli-Parisi[63] evolution equations. For example, for a non-singlet combination of fragmentation functions $D^{\pi^+} - D^{\pi^-}$:

$$Q^2\frac{\partial}{\partial Q^2} D^{\pi^+-\pi^-}(z,Q^2) = \frac{\alpha_s(Q^2)}{2\pi}\int_z^1 \frac{dy}{y} P_{q\to q}\left(\frac{z}{y}\right) D^{\pi^+-\pi^-}(z,Q^2) + O(\alpha_s(Q^2))^2 \tag{36a}$$

Since the splitting function $P_{q\to q}$ has the same form as in deep inelastic scattering, the moments of the fragmentation functions have similar logarithmic behaviours:

$$D_M^{\pi^+-\pi^-}(Q^2) = \int_0^1 dz\, z^{M-1} D^{\pi^+-\pi^-}(z,Q^2) \simeq \bar{D}_M \left(\ln Q^2/\Lambda^2\right)^{-d_M}\left[1+O(\alpha_s(Q^2))\right] \tag{36b}$$

where the anomalous dimensions d_M are the same as in deep inelastic scattering. Most of the $O[\alpha_s(Q^2)]$ radiative corrections in (34) have been calculated. The violation of factorization to which they lead[16] can best be expressed by taking double moments of the inclusive cross-sections (34) in both X and Z:

$$D_{MN}(Q^2) = \int_0^1 dx \int_0^1 dz\, x^{N-1} z^{M-1} \frac{d^2\sigma}{dx\,dz} \simeq \left(\ln Q^2/\Lambda^2\right)^{-d_M-d_N} \bar{D}_M \bar{D}_N \left[1 + C_{MN}\frac{\alpha_s(Q^2)}{2\pi} + \dots\right] \tag{37}$$

Figures 17a and 17b show respectively the total connection coefficient C_{MN} (37) and the extent to which factorization is broken (pieces of C_{MN} of the form a_M and b_N are subtracted). We see that a large breakdown of factorization is expected at moderate Q^2, even for relatively small values of M and N. Data from BEBC neutrino experiments[80] find qualitative agreement with perturbative QCD expectations already in the Q^2 range of 1 to 10 GeV^2.

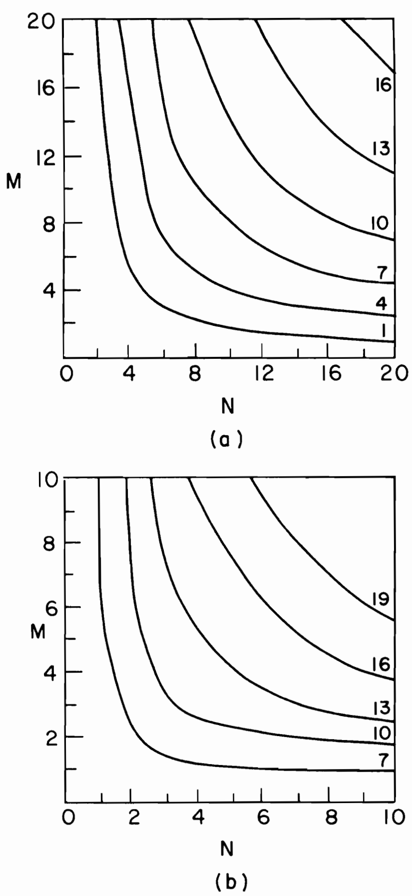


Fig. 17 The coefficients (37) of $\alpha_s(Q^2)/2\pi$ corrections in $\ell + h \to \ell + h' + X$: a) of the total correction, and b) of the contribution that violates factorization. Taken from Altarelli, R.K. Ellis, Martinelli and Pi (Ref. 16).

2.2 Lepton-pair Production

As was reported at this symposium[81], many of the qualitative features of the naive Drell-Yan $q\bar{q}$ annihilation mechanism[17] for lepton-pair production [$\sigma \propto A^1$, angular distribution $\alpha(1+\cos^2\theta)$, $\sigma(\pi^+N)/\sigma(\pi^-N) \to 1/4$, ...] are present in the experimental data. Perturbative QCD retains the bulk of these predictions[14], but has many sources of renormalization of the basic cross-section:

$$\begin{aligned}
M^2\frac{d^2\sigma}{dM\,dx_F} = {} & \frac{4\pi\alpha^2}{3.3.M^2}\frac{x_1x_2}{\sqrt{x_F^2+4x_1x_2}}\, q(x_1,Q^2)\bar{q}(x_2,Q^2) \times \left[1+\frac{\alpha_s(Q^2)}{2\pi}\frac{4}{3}\left(1+\frac{4\pi^2}{3}\right)\right] \\
& + \int_{x_1}^1 dy_1 \int_{x_2}^1 dy_2\, q(y_1,Q^2)\bar{q}(y_2,Q^2)\,\sigma(\alpha_s(Q^2)) \\
& + \int_{x_1}^1 dy_1 \int_{x_2}^1 dy_2 \left\{\left[q(y_1,Q^2)+\bar{q}(y_1,Q^2)\right]g(y_2,Q^2) + (y_1 \leftrightarrow y_2)\right\} \text{"}\sigma(\alpha_s(Q^2))\text{"} \\
& + O(\alpha_s(Q^2))^2
\end{aligned}$$

Notice in the first line of Eq. (38) the factor of 1/3 from colour, and the Q^2 dependence of the annihilating q and $\bar{q}$ distributions ($X_F = X_1 - X_2$ and $Q^2 = M^2 = sX_1X_2$). In the second line of Eq. (38) is the renormalization of the annihilation diagram due to vertex corrections and soft gluon corrections (see Fig. 18a). The portion $\propto \pi^2$ comes directly from the continuation from space-like to time-like Q^2 of $(\ln Q^2/p^2)^2$ terms involved in comparing deep inelastic scattering and the Drell-Yan process. This correction factor is very large[73]. For example, if we take $\alpha_s(Q^2) = 0.37$, a not unreasonable value in the presently accessible range of

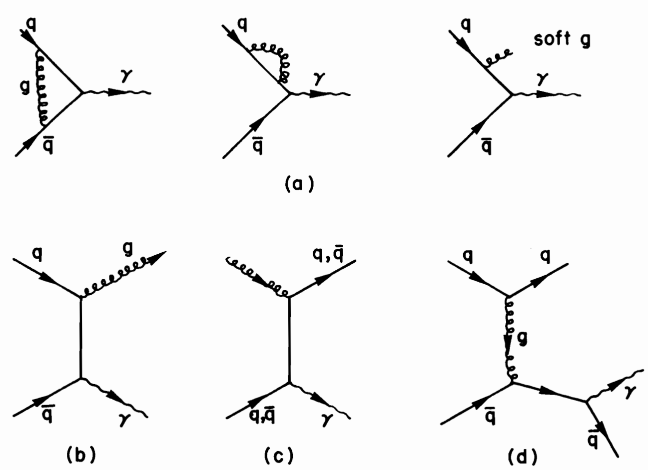


Fig. 18 Corrections of $O[\alpha_s(Q^2)]$ to the Drell-Yan cross-section from: a) vertex corrections and soft-gluon emission; b) hard-gluon emission; c) gq collisions; and d) qq collisions.

Q^2, then the correction is $1 + (1.1)$. Since the first-order radiative correction is so large at present Q^2, we cannot legitimately expect that the higher-order radiative corrections are negligible. It is clearly very important to try to get some handle on these higher-order terms. In particular, it may be that all or part of the large π^2 terms exponentiate, analogously to $(\ln Q^2/p_T^2)^2$ terms in the p_T distribution of Drell-Yan pairs[82]. [We may note facetiously that $e^{1 \cdot 1} = 3$, neatly cancelling the colour factor in the first line of Eq. (38)!]. The third-line terms in Eq. (38) refer to contributions with hard gluestrahlung (see Fig. 18b), where one or the other of the incoming q and $\bar{q}$ starts with a longitudinal momentum fraction $Y_i \neq X_i$. The magnitude of this correction relative to the lowest-order diagram depends on the M^2 and X_F of the observed $\ell^+\ell^-$ pair, as well as on the projectile and target used. Figure 19 shows that it can be important[83] in $pN \to \ell^+\ell^- + X$ at $\sqrt{s} = 27$ GeV, at $X_F = 0$ for large $\tau \equiv M^2/s$. The qg and $\bar{q}$g terms (Fig. 18c) are not so important; in fact Fig. 19 shows[83] that in pN collisions they are small and have a negative sign. The negative sign need not shock us since the "σ" in Eq. (38) has its (positive) leading logarithms removed and absorbed into the $q\bar{q}$ contribution. Adding together all the $O[\alpha_S(Q^2)]$ terms in Eq. (38), Fig. 19 shows that in pN collisions they are expected to give $\Delta\sigma/\sigma_0$ O(1) for all values of τ. A calculation has also been made[84] of some of the $O[\alpha_S^2(Q^2)]$ terms coming from qq scattering (see Fig. 18d). This calculation is not complete in the absence of calculations of higher-order radiative corrections to the fundamental subprocesses exhibited explicitly in Eq. (38). However, it indicates effects which may be small at presently accessible values of $\tau \lesssim 0.6$, but O(1) for τ near 1.

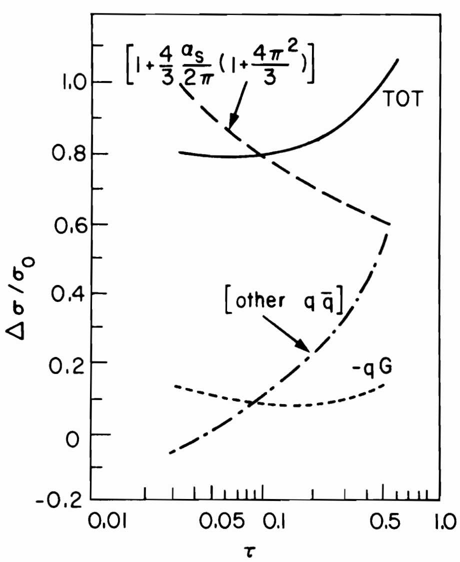


Fig. 19 Estimates[83] of $O[\alpha_S(Q^2)]$ corrections $\Delta\sigma$ to the naive Drell-Yan cross-section σ_0 for pN collisions at $\sqrt{s} = 27$ GeV.

Perturbative QCD also predicts the production of $\ell^+\ell^-$ pairs at large p_T accompanied by a large p_T gluon (Fig. 18b) or quark (Fig. 18c) jet[73]. We return later to attempts[82] to describe the distribution when $p_T \ll Q$. A suggestion going beyond conventional perturbative QCD is that higher-twist effects (see Fig. 20) may also make big corrections to the Drell-Yan cross-section at large values of X_F, particularly in πN collisions[85]. According to this analysis the π structure function (modulo QCD radiative corrections) should resemble

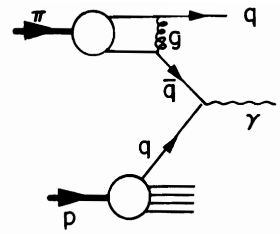


Fig. 20 A possibly[85] important higher-twist correction to $\pi N \to (\ell^+\ell^-) + X$.

$$(1-x)^2 + \frac{K^2}{Q^2} \tag{39}$$

Present data do not seem to be sensitive to the scaling violations generated by the $1/Q^2$ behaviour in formula (39)[85]. However, there is an associated expectation of a deviation of the $\ell^+\ell^-$ pair angular distribution from the naive $(1 + \cos^2\theta)$. One experiment[86] does see such a deviation from the naive angular distribution (see Fig. 21), but it is not yet clear whether this is indeed due to higher-twist effects or whether it could be an effect of higher orders in QCD perturbation theory.

It was reported at this conference[81] that several experiments see an apparent excess in the $\ell^+\ell^-$ pair cross-section by comparison with the naive $q\bar{q}$ pair annihilation mechanism. This is an O(1) effect comparable to the first-order perturbative QCD corrections shown in Fig. 19. Are the two related? It

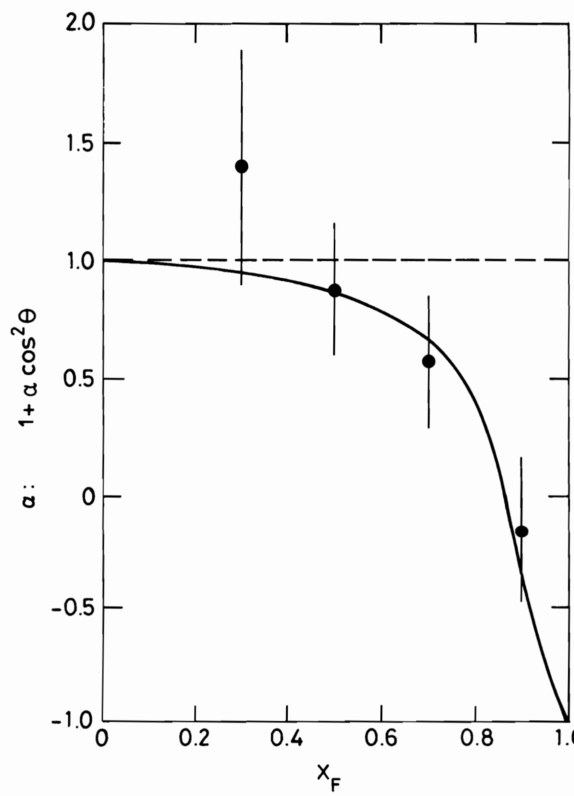


Fig. 21 The observed[86] angular dependence of $(\mu^+\mu^-)$ pairs in πN collisions, and the dependence estimated[89] from higher-twist effects, as a function of the longitudinal momentum fraction X_F of the $(\mu^+\mu^-)$ pair.

remains to be seen whether the experimental renormalization is a universal factor or whether it varies with X_F and Q^2 as expected by perturbative QCD. [Only the second-line correction term in Eq. (38) is universal, and it decreases $\propto 1/\ln Q^2$ at large Q^2.] There is also the important theoretical task of seeking a handle on the second- and higher-order QCD radiative corrections which are probably not negligible, and may render specious the present apparent consistency of theory and experiment. Life is certainly different from deep inelastic scattering, where the higher-order QCD corrections are rather difficult to pick out experimentally[51].

2.3 Jets

In any hard-scattering process the predominant QCD radiative corrections are those due to collinear gluestrahlung and pair creation (Fig. 5). These give rise to dominant configurations of jets of partons (Fig. 22)

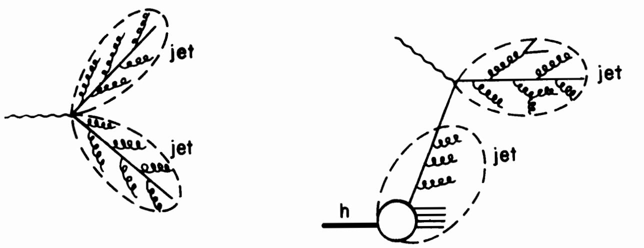


Fig. 22 Dominant two-jet configurations in $e^+e^- \to$ $\to$ hadrons and $\ell + h \to \ell' +$ hadrons.

-- two jets in e^+e^- annihilation, one forward and one backward jet in deep inelastic leptoproduction, and so on[68]. Wide-angle gluestrahlung and pair creation give extra jets (Fig. 23). If the jets are suitably defined, e.g. by an angular cut-off[18,19], then the production of each extra jet costs an extra factor of $\alpha_s(Q^2)$ so that, for example, in e^+e^- annihilation

$$\sigma(2\,\text{jet}) : \sigma(3\,\text{jet}) : \sigma(4\,\text{jet}) = 1 : \alpha_s : \alpha_s^2$$

-- a jet perturbation theory[18,68]. Explicitly, the three-jet cross-section is[18,37]

$$\frac{1}{\sigma_{tot}} \frac{d^2\sigma}{dX_q\, dX_{\bar q}} = \left(\frac{2\alpha_s}{3\pi}\right) \frac{X_q^2 + X_{\bar q}^2}{(1-X_q)(1-X_{\bar q})} + \cdots \qquad (40)$$

where $X_{q,\bar q} = 2E_{jet}/Q$. We will now study the phenomenology of wide-angle gluestrahlung, with particular reference to the exciting new PETRA data[87] reported[88] at this symposium.

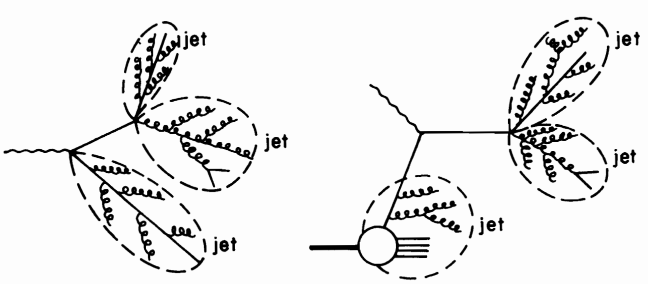


Fig. 23 Subdominant three-jet configurations in $e^+e^- \to$ $\to$ hadrons and $\ell + h \to \ell' +$ hadrons.

A first prediction is that there should be a large p_T cross-section in e^+e^- annihilation, where p_T is measured for example with respect to the thrust or sphericity axis. There should be scaling in the form

$$\frac{1}{\sigma_{tot}} \frac{d\sigma}{dp_T^2} = \frac{1}{Q^2} f(X_T) \times \left(\begin{matrix}\text{logarithmic}\\ \text{corrections}\end{matrix}\right) \qquad (41)$$

where $X_T = 2p_T/Q$. Indeed a big and increasing large p_T cross-section has been reported by the TASSO Collaboration[87,88]. Figure 24 shows their data, with an eyeball

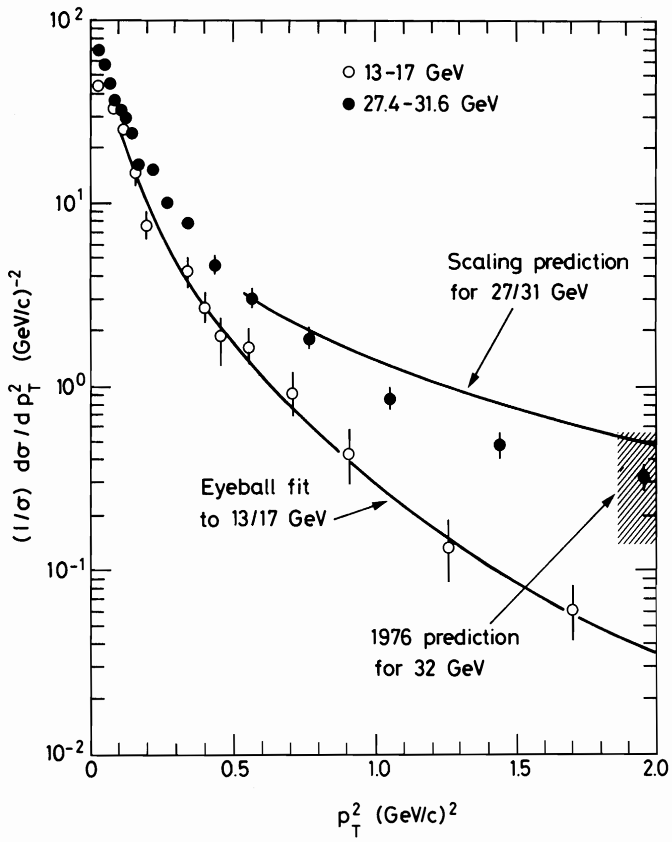


Fig. 24 The p_T distributions found by the TASSO Collaboration[87] at different centre-of-mass energies, compared with an eyeball fit to low-energy data, a related scaling prediction for high-energy data, and an interpolation of a 1976 prediction[18] based on $e^+e^- \to q\bar q g$.

fit to their low-energy (Q = 13 to 17 GeV) points compared with their high-energy (Q = 27.4 to 31.6 GeV) points. At $p_T > 1$ GeV, scaling is broken by (50 to 100)%, but this may be partly due to the logarithmic corrections expected in Eq. (41). Also shown for comparison is an interpolation from larger p_T of a 1976 absolute prediction[18] for the magnitude of the large p_T cross-section, based on the three-jet cross-section (40) and the assumption that quark and gluon jets fragment similarly -- an assumption supported by neutrino production[80] and T-decay data[71] at lower Q^2. Encouraged by this quantitative success of a QCD three-jet prediction, we are then led to ask whether the observed[87,88] structures (Fig. 25) are indeed the much-heralded QCD jets[68].

First off, the third QCD jet is due to a vector gluon, and we have no experimental evidence for this. There is some evidence for the vector nature of gluons

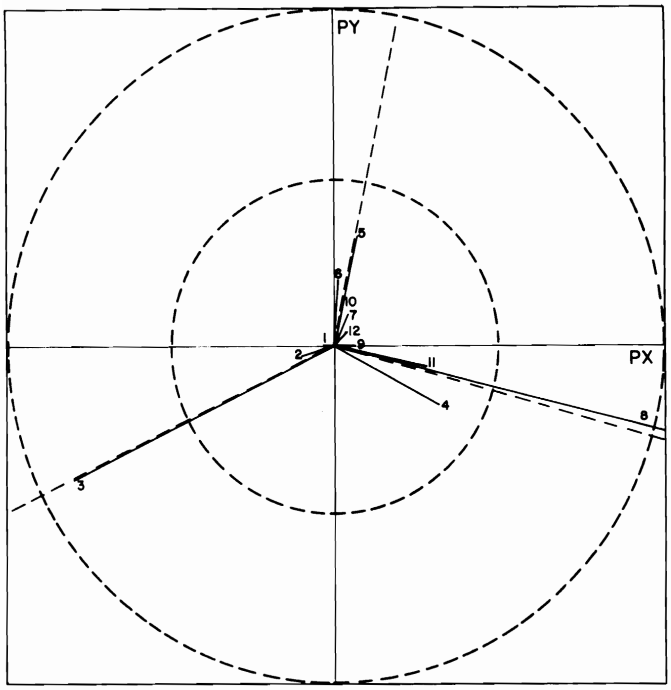


Fig. 25 Three-jet event reported by the TASSO Collaboration[87] at the Bergen Neutrino '79 conference in June 1979.

from deep inelastic scaling violations[8,42] and from T decay[71,89]: to test the vector nature in the $e^+e^- \to q\bar{q}g$ process, we should look at the angular distributions of three-jet events. For example, the normals to the planes of vector gluon three-jet events should[37] have a distribution

$$\frac{dN}{d(\cos\Theta)} \propto \left(2 + \sin^2\Theta\right) \tag{42}$$

We may also look at the angles in the planes. For example, if we look at three-jet events in the centre of mass of the two less energetic jets at their angular distribution relative to the axis of the most energetic jet, we find[90] the results shown in Fig. 26. Vector gluestrahlung gives a peaked distribution which is approximately

$$\frac{dN}{d(\cos\tilde{\Theta})} \propto \left(1 + 2\cos^2\tilde{\Theta}\right) \tag{43a}$$

whereas scalar gluestrahlung has approximately

$$\frac{dN}{d(\cos\tilde{\Theta})} \propto \left(1 + 0{\cdot}2\cos^2\tilde{\Theta}\right) \tag{43b}$$

and onium → 3-vector gluon jet decay has approximately

$$\frac{dN}{d(\cos\tilde{\Theta})} \propto \left(1 - 0{\cdot}1\cos^2\tilde{\Theta}\right) \tag{43c}$$

After testing the predictions (42) and (43) we will know whether vector gluons are being observed.

The next question is whether the gluon coupling is asymptotically free (1) as expected for QCD. To tell this really requires a large Q^2 range, and it is likely

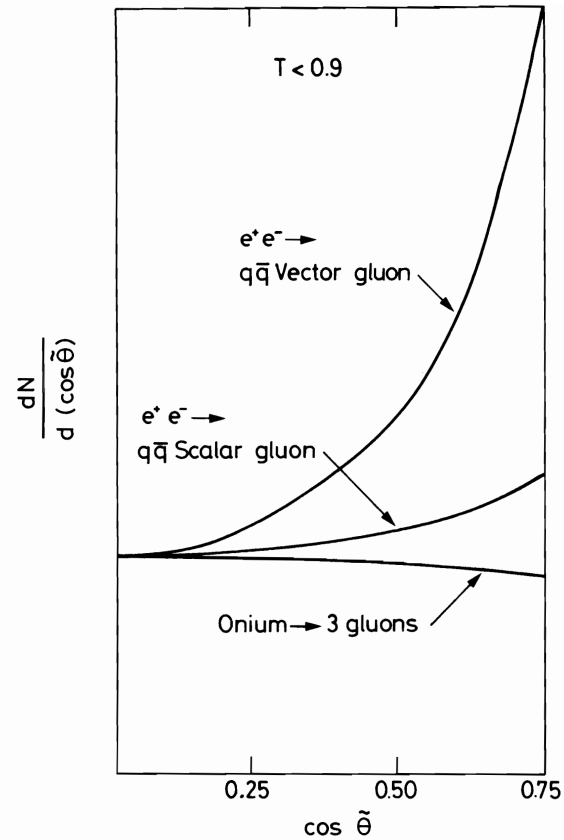


Fig. 26 Angular distributions in the centre-of-mass of the two least energetic jets expected[90] for $e^+e^- \to q\bar{q}$ + + vector gluon or scalar gluon, and onium → 3 gluons.

that the range obtainable at PETRA and PEP will not be sufficient. Perhaps LEP will give us a long enough handle to enable us to see a decrease in $\alpha_s(Q^2)$, but this is not clear either. We may have to content ourselves with the qualitative fact that $\alpha_s(Q^2)$ is small at large Q^2 so that perturbation theory is applicable, as suggested by the PETRA analyses[87,88].

A predicted aspect of QCD jets is the violation of scaling in the fragmentation functions $D(Z,Q^2)$ [79]. It is predicted that the scaling violations in gluon jets should be larger than those in quark jets[91], because the probability of collinear gluestrahlung from a gluon is larger than that from a quark. Observations of gluon jets at different Q^2 may reveal this phenomenon.

Another predicted aspect of QCD jets is that they should not have fixed p_T. For example, the cross-section $\sigma(\varepsilon,\delta,Q^2)$ for events in which less than a fraction ε of the total energy emerges outside two oppositely directed cones of opening angle δ is predicted[19] to obey

$$f(\varepsilon,\delta) \equiv \frac{\sigma(\varepsilon,\delta,Q^2)}{\sigma_{tot}} = 1 - O(\alpha_s(Q^2)) \approx 1 - O\left(\frac{1}{\ln Q^2}\right) \tag{44}$$

If we require that $f(\varepsilon,\delta)$ be constant, and keep ε also fixed, we find[92] that the cone opening angle

$$\delta(Q^2) \approx \left(\Lambda^2/Q^2\right)^{p(\varepsilon,\delta)} \tag{45}$$

where $p(\varepsilon,f)$ is a calculable power which is in general $< 1/2$, indicating that the typical p_T of hadrons grows with Q. Neither of the effects (44), (45) has yet been seen; present-day jets seem to have fixed p_T and are probably largely non-perturbative in origin[93]. This non-perturbative origin is perhaps reflected in the apparent similarity of gluon and quark jet widths as deduced from Υ decays[71], whereas perturbative QCD predicts that, asymptotically, gluon jets should be wider than quark jets[94].

Since none of the above QCD phenomena have been demonstrated at PETRA, it seems that QCD has not yet been proven by jets. It is not even clear that the existence of any type of gluon has been demonstrated. Even a simple uncorrelated jet model with a power-law p_T cut-off yields quite a few three-jet events[95]. Some people[96] have argued that the best way to test asymptotically free perturbation theory in e^+e^- annihilation may not be via jets at all, but just via measurements of the angular distribution of hadronic energy and of energy correlations. However, the recently observed three-jet events[87,88] certainly have plenty of dramatic value, and they may well turn out to have been the discovery of the QCD gluon.

2.4 Photons

As mentioned earlier, photons involved in 'hard' processes may either interact directly, or through point-like distribution and fragmentation functions, or through a 'soft' hadronic component. The distinction between the second and third classes was first seen in a renormalization group/operator product expansion analysis of $\gamma^*\gamma$ scattering[78]. These results have been reproduced and extended in a diagrammatic analysis of hard processes involving photons[14]. Ladder diagrams predominate: those where all loop momenta are in the $k^2 > \Lambda^2$ perturbative regime add up to the point-like forms (Figs. 16b, 16c)

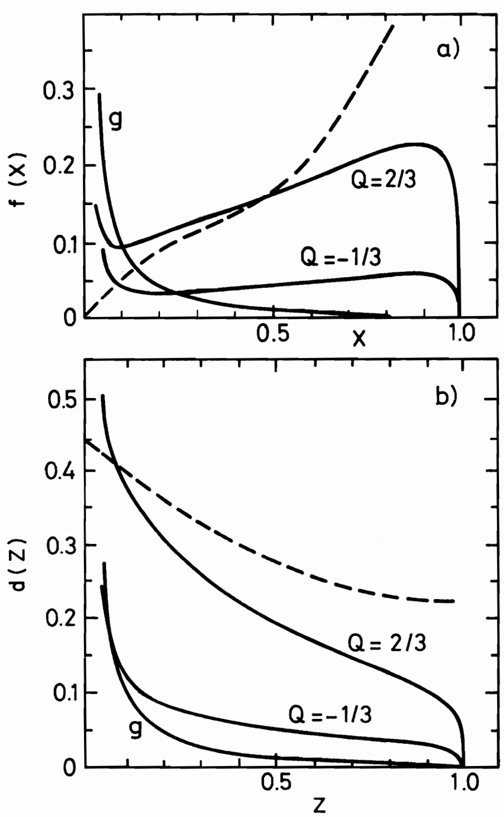


Fig. 27 a) Structure function for $\gamma \to q$, the dashed line corresponding to the absence of QCD corrections, and b) ditto for $q \to \gamma$ fragmentation functions[20].

$$\left.\begin{matrix} q_\gamma(X,Q^2) \\ D_{q\to\gamma}(Z,Q^2) \end{matrix}\right\} \underset{Q^2\to\infty}{\approx} \frac{\alpha_{e.m.}}{\pi}\ln Q^2 \left\{\begin{matrix} f(X) \\ d(Z) \end{matrix}\right\} + \cdots \tag{46}$$

Those where the nested-loop momenta k_i descend into the low k^2 region before reaching the photon, fall into the 'soft' hadronic part. In the language of the Altarelli-Parisi[63] evolution equations, the point-like piece is due to an electromagnetic driving term[97]

$$Q^2 \frac{\partial}{\partial Q^2} q_\gamma(X,Q^2) \approx e_q^2 \frac{\alpha_{e.m.}}{2\pi}\left(X^2+(1-X)^2\right) + \frac{\alpha_s(Q^2)}{2\pi}\int_X^1 \frac{dY}{Y} P_{q\to q}\left(\frac{X}{Y}\right) q_\gamma(Y,Q^2) + \cdots \tag{47}$$

Shown in Fig. 27 are calculations of the exactly computable 'reduced functions' f(X) and g(Z). For comparison, the dashed lines are the distributions as they would be in the absence of QCD renormalization (the gluon ladders in Fig. 16). The first-order QCD radiative corrections to $\gamma^*\gamma$ scattering have also been computed[98]. They turn out to be larger than those characteristic of deep inelastic scattering from a hadronic target, and particularly important at large X.

Among the most interesting applications of perturbative QCD to processes involving real photons are the production of two, three, and four jets in real $\gamma\gamma$ scattering[99] (see Fig. 28). The two-jet cross-section

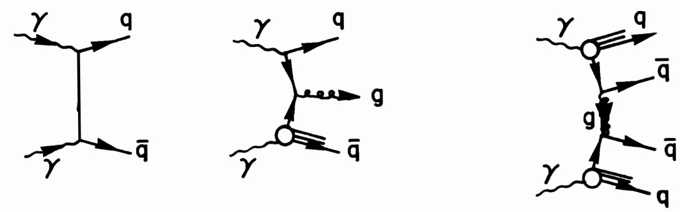


Fig. 28 Two-, three-, and four-jet contributions[99] to $\gamma\gamma \to$ hadrons.

scales proportionately to $\sigma(\gamma\gamma \to \mu^+\mu^-)$, the ratio just being

$$\frac{\sigma(\gamma\gamma \to 2\,\text{large } p_T \text{ jets})}{\sigma(\gamma\gamma \to \mu^+\mu^-)} \to 3 \sum_q e_q^4 \tag{48}$$

Perhaps surprisingly, the three-jet and four-jet cross-sections also scale in the same way -- there are no relative powers of $1/\ln Q^2$ as in $e^+e^- \to \gamma^* \to 3,4$ jets. The extra $1/\ln Q^2$ coming from the hard QCD vertex is cancelled by a $\ln Q^2$ coming from a 'soft' propagator.

At present, very little data exist on 'hard' processes involving photons, and it would be very interesting to look for some confirmation of the perturbative QCD predictions[100].

3. Exclusive Processes

One of the most exciting developments in recent months has been the growing realization[21,101] that perturbative QCD can be applied to a large number of exclusive processes at large momentum transfers -- examples are elastic form factors and wide-angle elastic scattering.

As far as pion form factors are concerned, it can be shown that in a light-like gauge the dominant

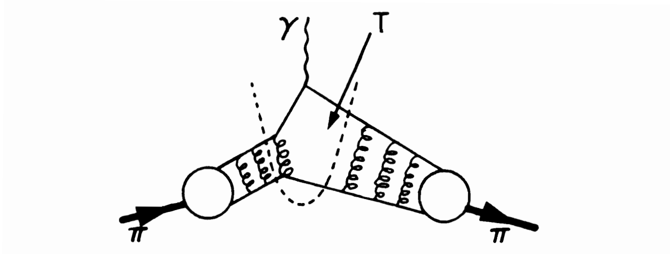


Fig. 29 Leading-order contribution to the meson form factor in QCD.

diagrams involve the two-constituent $q\bar{q}$ component of the wave function. The dominant diagrams at large momentum transfers are the generalized ladder diagrams of Fig. 29. Their first effect is to yield an evolution equation[102] for the meson wave function at large momenta:

$$X_1 X_2 \left\{ \frac{\partial}{\partial \xi} \phi(X_i, Q^2) + \frac{4}{33-2f} \phi(X_i, Q^2) \right\} = \int_0^1 dY_1 \int_0^1 dY_2\, \delta(1-Y_1-Y_2)\, V(X_i, Y_i)\, \phi(Y_i, Q^2) + \cdots \tag{49}$$

where X_1 and X_2 (Y_1 and Y_2) are fractions of longitudinal momentum carried by the quark and antiquark in the meson, $\xi = \ln\ln Q^2/\Lambda^2$, and $V(X_i,Y_i)$ represents the kernel due to one-gluon exchange. In the case of the pion, the vertex T of the virtual photon is quite simple (Fig. 29) and the form factor has the general structure

$$F_\pi(Q^2) \propto \phi^* T \phi \tag{50}$$

Asymptotically[21]

$$F_\pi(Q^2) \rightarrow \frac{16\pi\, \alpha_s(Q^2)\, f_\pi^2}{Q^2} \tag{51}$$

where f_π is the usual $\pi \rightarrow \mu\nu$ decay constant ($\simeq$ 93 MeV). Non-leading corrections to (51) can also be computed. They are controlled by the same anomalous dimensions of twist-2 $\bar{q}q$ operators that appear in deep inelastic scattering. Indeed, one can do the entire analysis of the π form factor using the operator product expansion[103], under suitable assumptions about the behaviour of the non-perturbative aspect of the pion wave function. In the case of the vector form factor of the π, all non-leading logarithms can be controlled using the renormalization group, but this may not be true for other form factors[104].

When we come to the nucleon form factors, the analysis proceeds analogously with only slightly more complications (see Fig. 30). The dominant component of the nucleon is the simplest qqq part[102]. It gives a

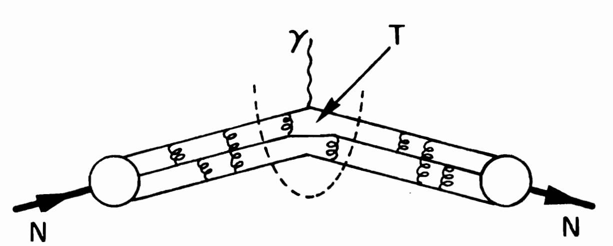


Fig. 30 Leading-order contribution to the baryon form factor in QCD.

power-law fall-off with Q^2 consistent with dimensional counting[105], modified by logarithmic factors which are slightly more complicated than in the π case, reflecting the greater complication of the wave function evolution equation and photon matrix element (see Fig. 30). The leading behaviour is found[102] to be

$$G_M(Q^2) \propto \frac{[\alpha_s(Q^2)]^2}{Q^4} \left(\ln Q^2/\Lambda^2\right)^{-\frac{4}{33-2f}} \tag{52}$$

There are still questions about the magnitude of the subasymptotic corrections to elastic form factors. Theoretically, they depend on unknown aspects of the hadron wave functions[21]. Experiments suggest they should be important at present Q^2, because the leading-order QCD predictions do not compare unequivocally well with the experimental data (Fig. 31). It would be nice

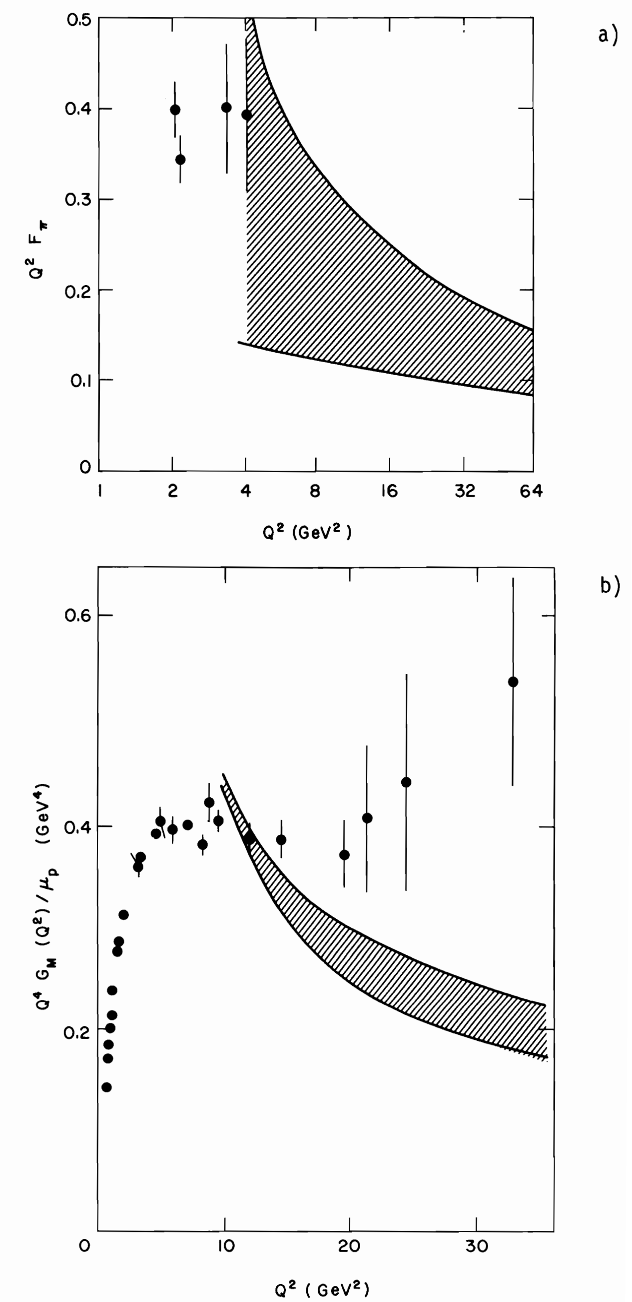


Fig. 31 QCD predictions[21] for a) the π form factor and b) the p form factor; the shaded areas representing subasymptotic uncertainties.

to have on the one hand more theoretical understanding of the subasymptotic corrections to the asymptotic results (51) and (52), and on the other hand to have better data at high Q^2 in either the space-like and/or the time-like region. Even if one cannot measure $e^+e^- \to \pi^+\pi^-$ at high time-like Q^2 because of the background from $e^+e^- \to \mu^+\mu^-$, perhaps one could measure $e^+e^- \to K^+K^-$ or $p\bar{p}$.

The analysis of elastic form factors can be extended to many other exclusive processes at large momentum transfers, such as quasi-elastic form factors, the production of individual hadrons, exclusive weak decays of particles containing heavy quarks, and wide-angle elastic scattering. In general, dimensional counting laws[105] will be reproduced for the powers of Q^2, with calculable corrections which are[102] for cross-sections:

$$(\ln Q^2)^{-n+2-\frac{8}{33-2f}n_B} \tag{53}$$

where n is the total number of interacting constituents and n_B is the number of external baryons. One phenomenon worth keeping an eye on is the polarization asymmetry in wide-angle elastic scattering, which would be small in the conventional perturbative QCD approximation of exchanging just vector gluons, but is experimentally measured to be large in wide-angle pp collisions at a beam energy of 11.75 GeV[106]. It has even been suggested[107] that this problem may be the Nemesis of perturbative QCD.

4. New Directions in QCD Perturbation Theory

All the applications of perturbative QCD that we have discussed up to now have involved sums of the type

$$\sum_{n=1}^{\infty} A_n (\alpha_s \ln Q^2)^n \left[1 + O\left(\frac{1}{\ln Q^2}\right)\right] \tag{53}$$

and are either rigorously known to be described by the renormalization group, or have been shown to behave in a similar way directly in perturbation theory. Can we do anything else? In particular can we sum series of the type

$$\sum_{n=1}^{\infty} B_n (\alpha_s \ln^2 Q^2)^n \left(1 + O\left(\frac{1}{\ln Q^2}\right)\right] \tag{54}$$

which appear in the study of various physically interesting phenomena?

An example where a summation of the type (54) has been achieved is the multiplicity of heavy quarks h in e^+e^- annihilation[108], which is presumably to be identified with the multiplicity of heavy mesons and baryons containing one of these heavy quarks h. The dominant contributions to the h multiplicity come from ladder diagrams, as in Fig. 32, where the final gluon

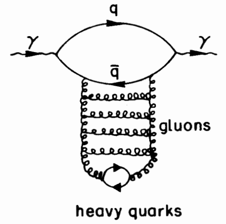


Fig. 32 Dominant[108] graph for the heavy hadron multiplicity in e^+e^- annihilation.

has an off-shell $Q_0^2 > 4m_h^2$. The summation of $(\alpha_s \ln^2 Q^2)^n$ terms (54) yields[108]

$$N_h(Q^2) \approx \left(\ln Q^2/Q_0^2\right)^{-1/4} \exp\sqrt{(\text{const}) \ln Q^2/Q_0^2} \tag{55}$$

This rises more slowly than any power of Q^2, but faster than any power of $\ln Q^2$, and is quite dramatic at high energies (Fig. 33). Such a rise would be dramatic evidence for the three-gluon vertex, which can be seen from Fig. 32 to play a vital role in generating the multiplicity curve. However, measuring the multiplicity of heavy quarks at enormous energies is a rather distant prospect, and although it is not (yet?) justified in QCD perturbation theory, one might try to apply formula (55) to ordinary hadrons by taking a cut-off $Q_0^2 = O(1)$ GeV^2, and multiplying the formula by some over-all factor to take into account the uncomputable hadronization of a low-mass gluonic cluster. Such a procedure gives a hadron multiplicity rising faster than the present data[88], and may indeed be incorrect.

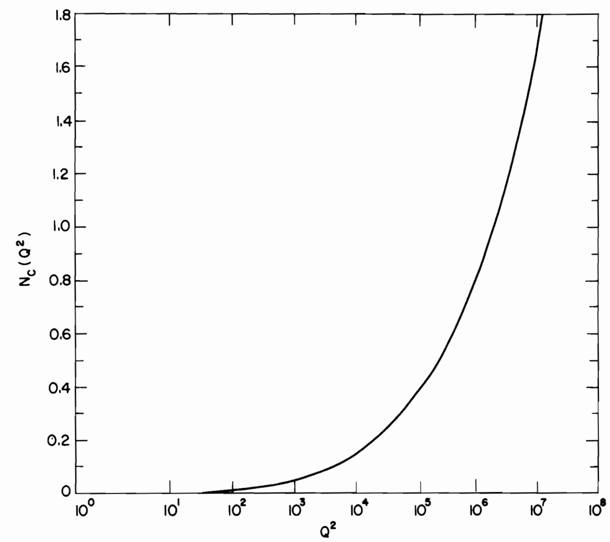


Fig. 33 Expected[108] Q^2 dependence of the charm multiplicity in e^+e^- annihilation.

Another instance where we may go beyond simple logarithm summation is in the study of small to moderate p_T in 'hard' processes such as Drell-Yan[82]. In the region $\Lambda^2 \ll p_T^2 \ll M^2_{\ell^+\ell^-}$ we can re-sum perturbation theory to obtain a Sudakov form factor. Looking at the calculation in impact parameter space, we see that the contribution of the large distance region is suppressed. It has been conjectured that this result may be extended to small $p_T^2 = O(\Lambda^2)$, and it has been found that the sensitivity to a finite non-zero primordial $\langle p_T \rangle$ goes away as the process gets harder. The intuition behind this result is illustrated in Fig. 34. In a

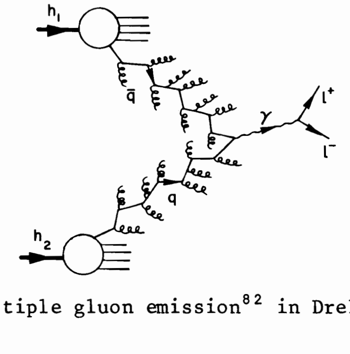
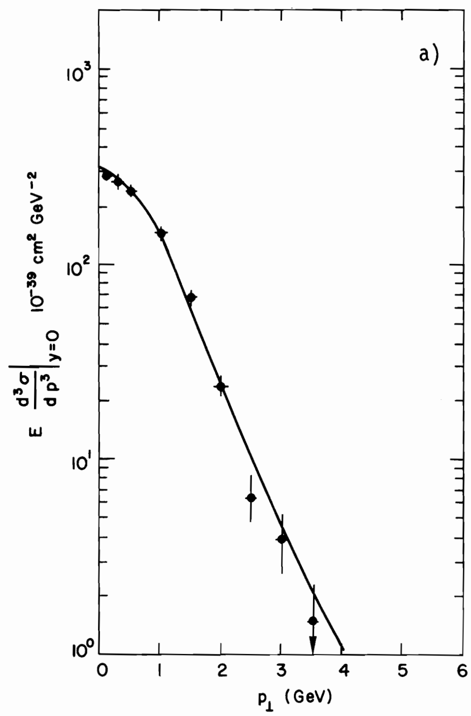


Fig. 34 Multiple gluon emission[82] in Drell-Yan collisions.

very hard Drell-Yan event, the incoming q and $\bar{q}$ will undergo multiple bremsstrahlung of semi-soft gluons. It might then be that the stochastic sum of the p_T engendered in each of these radiations dominates over the small initial $\langle p_T \rangle$ which gets 'lost in the crowd'. This possibility certainly fits the p_T distributions of Drell-Yan lepton pairs detected by the CFS (Fig. 35a)



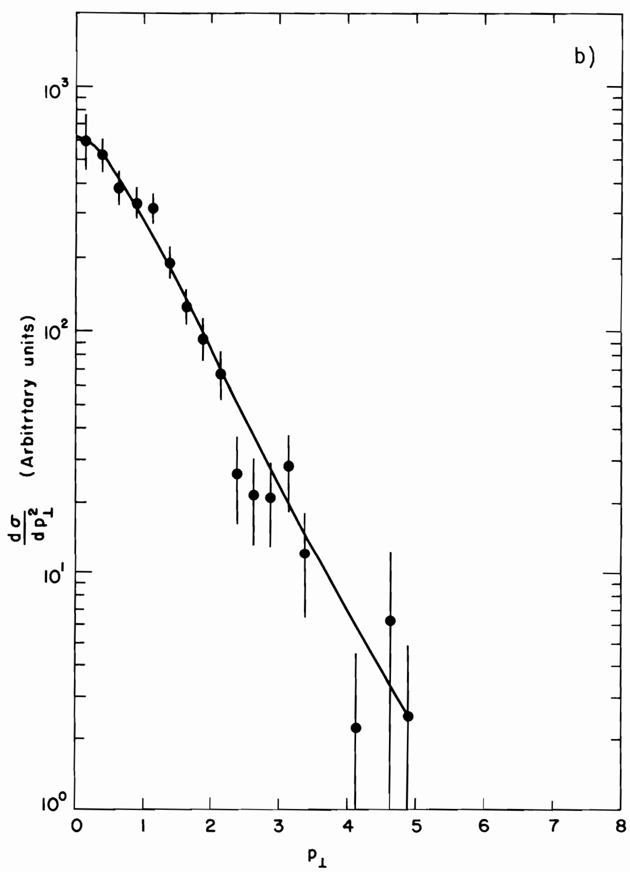


Fig. 35 QCD calculations[82,109] of the p_T distributions of Drell-Yan pairs: a) at $\sqrt{s}$ = 27 GeV, and b) at $\sqrt{s}$ = = 61 GeV, compared with experimental data.

and CHFMNP Collaborations[109] (Fig. 35b). The idea that even small p_T behaviour in 'hard' processes may be accessible to QCD perturbation theory is certainly very stimulating and worthy of further study.

Finally, the ambitious idea of "preconfinement"[110] should be mentioned. The suggestion is that perturbation theory may continue to generate quarks and gluons in a 'hard' process until such a stage that the entire hadronic final state can be covered by finite-mass colourless aggregates of quarks and gluons. These blobs would be "preconfined" in that non-perturbative effects would only have to operate over small momentum transfers in order to convert the perturbative final state into physical hadrons. This idea is very seductive, but the presently suggested way of realizing it is dependent on QCD containing quarks, and presumably a purely gluonic world should also "preconfine" and confine. Perhaps a more general realization of the "preconfinement" idea can be found.

5. Prospects and Problems

The theoretical status of perturbative QCD is very sound. The asymptotic behaviour at large momenta is well understood, in principle, for both inclusive[14] and at least a large class of exclusive[21] processes. Also, there is a systematic procedure for calculating subasymptotic corrections as a power series expansion in α_s, a quantity which vanishes as the momentum scale increases. A number of these subasymptotic corrections have been calculated, and they are of varying importance in different reactions, indicating that as one might have expected, the convergence of the perturbation series is non-uniform in Q^2. One theoretical shadow[69] on this rosy theoretical picture is that no general proof yet exists that non-perturbative effects do not modify the perturbative QCD predictions as is generally assumed.

Experimentally, there are various qualitative and even semi-quantitative pieces of evidence in favour of perturbative QCD, but as yet no convincing proof of its validity. Since the strong coupling α_s is so much larger than the weak and electromagnetic couplings at presently accessible values of Q^2, any approximation scheme is bound to converge much less rapidly than was the case for QED, and blockbustingly convincing tests analogous to (g-2) will be hard to find. Probably we have to resign ourselves to a long haul of piling up much circumstantial evidence in favour of QCD, rather than achieving swift conviction by finding a smoking gluon.

What are the immediate problems that seem interesting and important to investigate? Theoretically, it would be nice to see more results on

- radiative corrections to onium decays (e.g. $^3S_1 \to \to gg\gamma$, ggg; $^3P_{0,2} \to gg$; jet structures in final states);
- higher-twist effects in deep inelastic scattering (can one calculate them reliably, perhaps in a bag model?);
- beyond the next-to-leading order in Drell-Yan (can one get some control over the parts of these higher-order corrections which are not negligible?);
- higher-order effects in large p_T processes (are the corrections large as in Drell-Yan?);
- applications of the recent breakthrough in exclusive processes (e.g. to weak decays, elastic scattering);
- going beyond summing single logarithms $\Sigma_n (\alpha_s \ln Q^2)^n$ (e.g. for multiplicities, "preconfinement").

Experimentally there are many areas in which QCD perturbation theory can be subjected to significant tests. To name but a few:

- in e^+e^- annihilation: R, scaling violations in final states, multiple jet studies, onium decays;

- in deep inelastic scattering: more precise data at high and low Q^2 to try to separate QCD logarithms from higher-twist effects, σ_L/σ_T, final states;

- in Drell-Yan processes: the normalization question -- is the cross-section really larger than the naive $q\bar{q}$ annihilation model? and does any enhancement factor vary with M^2, beam type, X_F? angular distributions and scaling properties;

- studies of hard processes involving photons (e.g. photoproduction at large p_T, deep inelastic Compton scattering, final-state photons in deep inelastic processes, $\gamma\gamma$ collisions);

- exclusive processes (are there logarithmic deviations from the normal dimensional counting rules?);

In assessing the progress made and work ahead we should remember that so far there is precious little direct evidence for fundamental aspects of QCD, such as

- the vector nature of the gluons (some evidence comes from deep inelastic scaling violations[42] and from Υ decays[71], but none yet comes from the $e^+e^- \to 3$ jet analyses);

- asymptotic freedom (there is plenty of circumstantial evidence that α_s is small at $Q^2 > 1$ or 2 GeV^2, but only limited quantitative information from deep inelastic scaling violations and quarkonium studies);

- the three-gluon vertex which underlies asymptotic freedom and reflects the gauge nature of QCD. Possible ways to see its effects include scaling violations in gluon jets which should be larger than those in quark jets[91] [some tentative evidence from J/ψ and Υ decays[71]], scaling violations in the gluon distribution inside the nucleon [some evidence from deep inelastic scaling violations[9,10] -- another place to look would be σ_L/σ_T [111]], the width of gluon jets which should be broader than quark jets at large Q^2 [92], the multiplicities of heavy quarks[108,112], and asymmetries in heavy quarkonium decay[113].

We have every reason to hope that progress on these theoretical and experimental fronts will be rapid in the next two years, and that at the next lepton-photon symposium the rapporteur on QCD will be able to agree with the Chicago Tribune[114] that QCD is established.

Acknowledgements

It is a pleasure to thank R.M. Barnett, S.J. Brodsky, A. De Rújula, M.K. Gaillard, C.H. Llewellyn Smith, A. Para, D.H. Perkins, R. Petronzio, G. Preparata, C.T. Sachrajda and W.G. Scott for discussions about the material presented here.

References

1) For reviews, see
H.D. Politzer, Phys. Rep. 14C, 129 (1974).
W. Marciano and H. Pagels, Phys. Rep. 36C, 137 (1978).
2) J.D. Bjorken and E.A. Paschos, Phys. Rev. 158, 175 (1969).
S.D. Drell, D.J. Levy and T.-M. Yan, Phys. Rev. D 1, 1035 (1970).
R.P. Feynman, Photon-hadron interactions (Benjamin, New York, 1972).
3) See, for example, the reports by O. Nachtmann (p. 811), L. Susskind (p. 895) and S. Weinberg (p. 619), Proc. 1977 Int. Symp. on Lepton and Photon Interactions at High Energies, Hamburg, 1977, ed. F. Gutbrod (DESY, Hamburg, 1977).
4) A.M. Polyakov, Proc. Int. Symp. on Lepton and Photon Interactions at High Energies, Stanford, 1975, ed. W.T. Kirk (SLAC, Stanford, 1975), p. 855.
5) C.H. Llewellyn Smith, Proc. Int. Symp. on Lepton and Photon Interactions at High Energies, Stanford, 1975, ed. W.T. Kirk (SLAC, Stanford, 1975), p. 709.
6) H.L. Anderson, H.S. Matis and L.C. Myrianthopoulos, Phys. Rev. Lett. 40, 1061 (1978).
7) L. Hand, Proc. Int. Symp. on Lepton and Photon Interactions at High Energies, Hamburg, 1977, ed. F. Gutbrod (DESY, Hamburg, 1977), p. 417; and O. Nachtmann (Ref. 3) review data prior to the previous symposium, and make comparisons with QCD predictions.
8) P.C. Bosetti et al., Nucl. Phys. B142, 1 (1978).
9) J.G.H. de Groot et al., Z. Phys. C1, 143 (1979); Phys. Lett. 82B, 292, 456 (1979).
10) H.L. Anderson et al., Fermilab preprint Pub-79/30-EXP (1979).
11) T. Appelquist and H. Georgi, Phys. Rev. D 8, 4000 (1973).
A. Zee, Phys. Rev. D 8, 4038 (1973).
12) T. Appelquist and H.D. Politzer, Phys. Rev. Lett. 34, 43 (1975); Phys. Rev. D 12, 1404 (1975).
13) D.J. Gross and F. Wilczek, Phys. Rev. D 8, 3633 (1973) and D 9, 980 (1974).
H. Georgi and H.D. Politzer, Phys. Rev. D 9, 416 (1974).
14) D. Amati, R. Petronzio and G. Veneziano, Nucl. Phys. B140, 54 (1978) and B146, 29 (1978).
Yu.L. Dokshitzer, D.I. D'yakonov and S.I. Troyan, Proc. 13th Winter School of the Leningrad B.P. Konstantinov Institute of Nuclear Physics [available as SLAC translation No. 183 (1978)].
A.V. Efremov and A.V. Radyushkin, Serpukhov preprint E2-11535 (1978).
R.K. Ellis, H. Georgi, M. Machacek, H.D. Politzer and G.G. Ross, Phys. Lett. 78B 281 (1978); Nucl. Phys. B152, 285 (1979).
W.B. Frazer and J.F. Gunion, Phys. Rev. D 19, 2447 (1979).
Y. Kazama and Y.-P. Yao, Phys. Rev. Lett. 41, 611 (1978); Phys. Rev. D 19, 3111, 3121 (1979).
S. Libby and G. Sterman, Phys. Lett. 78B, 618 (1978); Phys. Rev. D 18, 3252, 4737 (1978).
C.H. Llewellyn Smith, Acta Phys. Austriaca Suppl. XIX, 331 (1978).
G. Sterman, Phys. Rev. D 19, 3135 (1979).
For a review, see
Yu.L. Dokshitzer, D.I. D'yakonov and S.I. Troyan, Hard processes in QCD, Phys. Rep. C (in press, 1979).
15) For a different formulation using cut vertices, see
A.H. Mueller, Phys. Rev. D 18, 3705 (1978).
Some applications of this formalism are described in
S. Gupta and A.H. Mueller, Phys. Rev. D 20, 118 (1979).

16) N. Sakai, Phys. Lett. 85B, 67 (1979).
R. Baier and K. Fey, Bielefeld Univ. preprint BI-TP 79/11 (1979).
G. Altarelli, R.K. Ellis, G. Martinelli and S.-Y. Pi, MIT preprint CTP 793 (1979).
17) S.D. Drell and T.-M. Yan, Phys. Rev. Lett. 25, 316 (1970).
18) J. Ellis, M.K. Gaillard and G.G. Ross, Nucl. Phys. B111, 253 (1976).
19) G. Sterman and S. Weinberg, Phys. Rev. Lett. 39, 1436 (1977).
20) E. Witten, Nucl. Phys. B120, 189 (1977).
S.J. Brodsky, T. DeGrand, J.F. Gunion and J.H. Weis, Phys. Rev. Lett. 41, 672 (1978); Phys. Rev. D 19, 1418 (1979).
W.R. Frazer and J.F. Gunion, Phys. Rev. D 20, 147 (1979).
C.H. Llewellyn Smith, Phys. Lett. 79B, 83 (1978).
T.A. DeGrand, Phys. Lett. 84B, 478 (1979).
21) A.V. Efremov and A.V. Radyushkin, Ref. 14.
G.L. Farrar and D.R. Jackson, Phys. Rev. Lett. 43, 246 (1979) report some partial results. A complete analysis and justification, including baryon form factors, is given by
S.J. Brodsky and G.P. Lepage, SLAC preprint PUB-2794 (1979); Phys. Rev. Lett. 43, 545 (1979) and SLAC preprint PUB-2343 (1979). An operator product expansion analysis given by
S.J. Brodsky, Y. Frishman, G.P. Lepage and C.T. Sachrajda (preprint TH.2731-CERN, 1979).
22) E.C.G. Stueckelberg and A. Peterman, Helv. Phys. Acta 26, 499 (1953).
M. Gell-Mann and F.E. Low, Phys. Rev. 95, 1300 (1954).
For reviews of the applications of the renormalization group to deep inelastic scattering, see
A. Peterman, Phys. Rep. 53C, 157 (1979).
A.J. Buras, Fermilab preprint PUB-79/17-THY (1979).
23) S. Coleman and D.J. Gross, Phys. Rev. Lett. 31, 851 (1973).
24) G. 't Hooft, unpublished (1972).
H.D. Politzer, Phys. Rev. Lett. 30, 1346 (1973).
D.J. Gross and F. Wilczek, Phys. Rev. Lett. 30, 1343 (1973).
25) K.G. Wilson, Phys. Rev. 179, 1499 (1969).
Y. Frishman, Ann. Phys. (N.Y.) 66, 373 (1971).
R.A. Brandt and G. Preparata, Nucl. Phys. B27, 541 (1971).
26) W.E. Caswell, Phys. Rev. Lett. 33, 244 (1974).
D.R.T. Jones, Nucl. Phys. B75, 531 (1974).
27) S.L. Adler, Phys. Rev. D 10, 3714 (1974).
A. De Rújula and H. Georgi, Phys. Rev. D 13, 1296 (1976).
E.C. Poggio, H.R. Quinn and S. Weinberg, Phys. Rev. D 13, 1958 (1976).
R.G. Moorhouse, M.R. Pennington and G.G. Ross, Nucl. Phys. B124, 285 (1977).
R. Shankar, Phys. Rev. D 15, 755 (1978).
28) M. Dine and J. Sapirstein, Phys. Rev. Lett. 43, 668 (1979).
K.G. Chetyrkin, A.L. Kataev and F.V. Tkachov, Phys. Lett. 85B, 277 (1979).
D.A. Ross, A. Terrano and S. Wolfram, private communication (1979).
29) G. 't Hooft and M. Veltman, Nucl. Phys. B44, 189 (1972).
30) W.A. Bardeen, A.J. Buras, D.W. Duke and T. Muta, Phys. Rev. D 18, 3998 (1978).
31) M. Dine and J. Sapirstein, Ref. 28; and R. Schwitters and J. Siegrist quoted therein.
32) R. Barbieri, E. d'Emilio, G. Curci and E. Remiddi, Nucl. Phys. B154, 535 (1979).
33) R. Barbieri, R. Gatto and E. Remiddi, Phys. Lett. 61B, 465 (1976).
34) SLAC-LBL Mark I and Mark II Collaborations, private communications (1979).
35) E. Bloom, Talk on results of the Crystal Ball experiment, these proceedings (1979).
36) For the indirect method of extracting the $^3P_0$ and $^3P_2$ widths, see
M.S. Chanowitz and F. Gilman, Phys. Lett. 63B, 178 (1976).
37) T. DeGrand, Y.J. Ng and S.-H.H. Tye, Phys. Rev. D 16, 3251 (1977).
M. Krammer and H. Krasemann, Phys. Lett. 73B, 58 (1978).
A. De Rújula, J. Ellis, E.G. Floratos and M.K. Gaillard, Nucl. Phys. B138, 387 (1978).
K. Koller and T.F. Walsh, Phys. Lett. 72B, 227 (1977) and 73B, 504 (1978).
S.J. Brodsky, D.G. Coyne, T.A. DeGrand and R.R. Horgan, Phys. Lett. 73B, 203 (1978).
H. Fritzsch and K.-H. Streng, Phys. Lett. 74B, 90 (1978).
K. Hagiwara, Nucl. Phys. B137, 164 (1978).
K. Koller, H. Krasemann and T.F. Walsh, Z. Phys. C1, 71 (1979).
For deviations from two- and three-jet structure in the e^+e^- continuum, see
A. Ali et al., Phys. Lett. 82B, 285 (1979).
38) N. Christ, B. Hasslacher and A.H. Mueller, Phys. Rev. D 6, 3543 (1972).
39) O. Nachtmann, Nucl. Phys. B63, 237 (1973) and B78, 455 (1974).
S. Wandzura, Nucl. Phys. B112, 412 (1977).
S. Matsuda and T. Uematsu, Kyoto Univ. preprint RIFP-376 (1979).
40) For a more leisurely discussion of this and other deep inelastic questions, see
J. Ellis, Invited talk at Neutrino 79, Int. Conf. on Neutrinos, Weak Interactions and Cosmology, Bergen, preprint TH.2701-CERN (1979).
41) G. Parisi, Phys. Lett. 43B, 207 (1973).
D.J. Gross, Phys. Rev. Lett. 32, 1071 (1974).
A.J. Buras and K.J.F. Gaemers, Nucl. Phys. B132, 249 (1978).
42) J. Ellis, Current trends in the theory of fields, eds. J.E. Lannutti and P.K. Williams (AIP, New York, 1978), p. 81; and Proc. 1978 SLAC Summer Institute on Particle Physics, SLAC-215 (1978).
See also M. Glück and E. Reya, Phys. Rev. D 16, 3242 (1977).
43) D. Bailin and A. Love, Nucl. Phys. B75, 159 (1974).
M. Glück and E. Reya, Ref. 42 and DESY preprint 79/13 (1979).
44) E.G. Floratos, D.A. Ross and C.T. Sachrajda, Nucl. Phys. B129, 66 (1977) and 139, 545 (1978).
W.A. Bardeen, A.J. Buras, D.W. Duke and T. Muta, Ref. 30.
E.G. Floratos, D.A. Ross and C.T. Sachrajda, Phys. Lett. 80B, 269 (1979); Nucl. Phys. B152, 493 (1979).
D.A. Ross and C.T. Sachrajda, Nucl. Phys. 149, 497 (1979).
W.A. Bardeen and A.J. Buras, Fermilab preprint Pub-79/31-THY (1979).
J. Kodaira, S. Matsuda, T. Muta, T. Uematsu and S. Sasaki, Phys. Rev. D 20, 627 (1979).
J. Kodaira, S. Matsuda, K. Sasaki and T. Uematsu, Kyoto preprint RIFP-360 (1979).
For reviews, see A. Peterman and A.J. Buras, Ref. 22.
For a first attempt at guessing the effects of still higher orders, see
M. Moshe, Tel-Aviv Univ. preprint TAUP-769-79 (1979).
45) L.F. Abbott and R.M. Barnett, SLAC preprint PUB-2325 (1979).
46) See, however,
A. De Rújula, H. Georgi and H.D. Politzer, Ann. Phys. (N.Y.) 103, 315 (1977).

47) M.R. Pennington and G.G. Ross, Oxford Univ. preprint 23/79 (1979).
For an earlier discussion, see
R. Barbieri, L. Caneschi, G. Curci and E. d'Emilio, Phys. Lett. 81B, 207 (1979).

48) For a qualitative discussion of these effects, see Ref. 45.

49) The data in this figure come via R.M. Barnett, Proc. 1979 SLAC Summer Institute on Particle Physics and SLAC-PUB-2396, 1979.

50) M. Bacé, Phys. Lett. 78B, 132 (1978); and for somewhat different definitions of Λ_N, see Refs. 10 and 30 and
D.W. Duke and R.G. Roberts, Phys. Lett. 85B, 289 (1979).

51) A. Para and C.T. Sachrajda, preprint TH.2702-CERN (1979).

52) The data indicated by crosses come from:
T.W. Quirk, J.B. McAllister, W.S.C. Williams and C. Tao, Oxford Univ. preprint, Non-singlet moments in charged lepton scattering and QCD (1979); while those indicated by circles are a guess by A. Para and the author of the likely effect of including the ln ln Q^2/Λ_N^2 dependence (19) into the parametrization. The other data points come from a re-analysis by A. Para of the data of Refs. 8 and 9 using uniform assumptions. For discussions of the effects of differences in the assumptions used in previous analyses (Refs. 8,9) see J. Wotschak, Talk given in Parallel Session 2 at the EPS Int. Conf. on High-Energy Physics, Geneva, 1979, and the invited talks by A. Para and W.S.C. Williams at this symposium. I am indebted to A. Para for many discussions on these points.

53) Note that D.W. Duke and R.G. Roberts (Ref. 50) used a different definition of Λ_N and that the N dependence they quote for their definition is incorrect (C.T. Sachrajda, private communication). Their extraction of the Λ_N from e and μ scattering data differs from that of T.W. Quirk et al. (Ref. 52) which is used here.

54) It would probably be better to use a value of f which is Q^2 dependent, as four quarks are surely not relevant to the low Q^2 data, whereas they should be relevant to high Q^2 data. Note also that the CDHS moment analysis and extraction of Λ_N combined F_2^{eN} and $F_3^{\nu N}$ scattering data, while QCD predicts (Ref. 51) that these should have different values of Λ_N.

55) M. Glück and E. Reya, DESY preprint 79/13 (1979).

56) D.H. Perkins, private communication (1979).

57) See also the analyses of Refs. 6, 8, 9, and 10.

58) P. Castorina, G. Nardulli and G. Preparata, preprint TH.2670-CERN (1979), make a similar point using a different form of subasymptotic corrections motivated by the massive quark model.

59) These fits are also equally good for the experimental values of $d_N/d_{N'}$: see Refs. 45 and 58 for details.

60) For a light-hearted but serious discussion of this and other deep inelastic issues, see:
A. De Rújula, J. Ellis, R. Petronzio, G. Preparata and W. Scott, Can one tell QCD from a hole in the ground? Play presented at the 1979 "Ettore Majorana" Int. School of Subnuclear Physics, Erice, 1979, CERN TH preprint in preparation (1979).

61) L. Baulieu and C. Kounnas, Nucl. Phys. B155, 429 (1979).

62) A. González-Arroyo, C. López and F.J. Ynduráin, Nucl. Phys. B153, 161 (1979); Madrid preprint FTUAM/79-3 (1979), and preprint TH.2728-CERN (1979). See also
D.A. Ross, CalTech preprint CALT-68-699 (1979) and R.D. Field, private communication.

63) G. Altarelli and G. Parisi, Nucl. Phys. B126, 298 (1977).

64) T. Goldman and D.A. Ross, Phys. Lett. 84B, 208 (1979).
W.J. Marciano, Phys. Rev. D 20, 274 (1979).

65) This class of applications of the parton model has been vigorously criticized by
G. Preparata, Lepton-hadron structure, ed. A. Zichichi (Academic Press, N.Y., 1975), p. 54.

66) The results one gets from leading-order QCD for large p_T processes [see, for example, the work of the CalTech group reviewed in R.D. Field, CalTech preprint CALT-68-696 (1978)] are often very sensitive to the choice of large kinematic invariant ($Q^2 = \hat{t}$, $\hat{s}$, $\hat{u}$, or some combination?) used. These choices differ only in non-leading logarithms, and so indicate that subasymptotic corrections to large p_T processes may be very large at moderate p_T.

67) See, for example,
I.G. Halliday, Nucl. Phys. B103, 343 (1976).
H.D. Politzer, Phys. Lett. 70B, 430 (1977) and Nucl. Phys. B129, 301 (1977).
C.T. Sachrajda, Phys. Lett. 73B, 185 (1978) and 76B, 100 (1978).
W. Furmański, Phys. Lett. 77B, 312 (1978).

68) See, in particular, the analyses by Llewellyn Smith, Frazer and Gunion, and Kazama and Yao in Ref. 14.

69) Non-perturbative instanton effects have been shown to be negligible in high-energy e^+e^- annihilations:
N. Andrei and D.J. Gross, Phys. Rev. D 18, 468 (1978).
T. Appelquist and R. Shankar, Phys. Lett. 78B, 468 (1978).
L. Baulieu, J. Ellis, M.K. Gaillard and W.J. Zakrzewski, Phys. Lett. 77B, 290 (1978).
R.D. Carlitz and C. Lee, Phys. Rev. D 17, 3238 (1978).
S. Shei, Nucl. Phys. B156, 225 (1979); and not to affect the leading scaling deviations expected for deep inelastic scattering:
L. Baulieu, J. Ellis, M.K. Gaillard and W.J. Zakrzewski, Phys. Lett. 81B, 41 (1979).
C.A. Flory, LBL Berkeley preprint 8499 (1979); but they may affect factorization in hard scattering processes:
J. Ellis, M.K. Gaillard and W.J. Zakrzewski, Phys. Lett. 81B, 224 (1979). Effects due to spectator interactions have been investigated:
A.V. Smilga, ITEP Moscow preprint ITEP-20 (1979).
C.T. Sachrajda and S. Yankielowicz (to be published) reported in
J. Ellis and C.T. Sachrajda, QCD and its applications, 1979 Cargèse Summer Institute Lectures (1979).

70) Slightly different definitions are adopted by Baier and Fey and by Altarelli, R.K. Ellis, Martinelli and Pi (Ref. 16). The latter authors point to an interesting partial verification of the reciprocity relation [V. Gribov and L. Lipatov, Sov. J. Nucl. Phys. 15, 675 (1972)] between deep inelastic scattering and e^+e^- annihilation. See, however,
L. Baulieu, E.G. Floratos and C. Kounnas, Ecole Normale Supérieure preprint LPTENS 79/16 (1979).

71) See the PLUTO analysis presented by
H. Meyer, Proc. Int. Symp. on Lepton and Photon Interactions at High Energies (FNAL, Batavia, 1979). The interest of these moments has been emphasized by
K. Koller, T.F. Walsh and P.M. Zerwas, Phys. Lett. 82B, 263 (1979).

72) E.G. Floratos, Nuovo Cimento 143, 241 (1978).
H. Georgi and H.D. Politzer, Phys. Rev. Lett. 40, 3 (1978).
A. Méndez, Nucl. Phys. B145, 199 (1978).
A. Méndez, A. Raychaudhuri and V.J. Stenger, Nucl. Phys. B148, 499 (1979).
P. Binétruy and G. Girardi, Nucl. Phys. B155, 150 (1979).
R.D. Peccei and R. Rückl, Phys. Lett. 84B, 95 (1979).
E.K. Manesis and N.A. Papadopoulos, preprint TH.2708-CERN (1979).
K.-H. Streng, T.F. Walsh and P.M. Zerwas, DESY preprint 79-10 (1979).
H. Georgi and J. Sheiman, Phys. Rev. D 20, 111 (1979).
F. Hayot and A. Morel, CEN-Saclay preprint DPh-T/111/79 (1979).
Chang C.-h., preprint TH.2736-CERN (1979).
For the first step towards a theoretical justification for the phenomenological processes used here, see
K.H. Craig and C.H. Llewellyn Smith, Phys. Lett. 72B, 349 (1978).
73) J. Abad and B. Humpert, Phys. Lett. 77B, 105 (1978) and 80B, 286 (1979).
G. Altarelli, R.K. Ellis and G. Martinelli, Nucl. Phys. B143, 521 (1978), erratum Nucl. Phys. B146, 544 (1978), and MIT preprint CTP-776 (1979).
A.P. Contogouris and J. Kripfganz, Phys. Rev. D 19, 2207 (1979).
K. Harada, T. Kaneko and N. Sakai, Nucl. Phys. B155, 169 (1979) and erratum to be published.
B. Humpert and W.L. van Neerven, Phys. Lett. 84B, 327 (1979), 85B, 293 (1979), and preprint TH.2738-CERN (1979) and Erratum.
J. Kubar-André and F.E. Paige, Phys. Rev. D 19, 221 (1978).
See the penultimate paper for a clarification of some conflicts in the previous literature.
74) It may be that the large primordial $\langle p_T \rangle$ found necessary by the CalTech group [see R.P. Feynman, R.D. Field and G.C. Fox, Nucl. Phys. B128, 1 (1977), and R.D. Field (Ref. 66)] may in part reflect non-leading perturbative effects: see the discussion of Drell-Yan p_T in Section 5 and R. Stroynowski, private communication and Proc. 1979 SLAC Summer Institute (in preparation) (1979).
75) Z. Kunszt and E. Pietarinen, DESY preprint 79/34 (1979).
T. Gottschalk, E. Monsay and D. Sivers, Argonne preprint ANL-HEP-79-66 (1979), and references therein.
76) J.D. Bjorken and E.A. Paschos, Ref. 2.
77) H. Fritzsch and P. Minkowski, Phys. Lett. 69B, 316 (1977).
78) See E. Witten, Ref. 20.
79) J.F. Owens, Phys. Lett. 76B, 85 (1978).
T. Uematsu, Phys. Lett. 79B, 97 (1978).
80) See D.H. Perkins, W.G. Scott and W. Wittek, Invited talks at the Neutrino 79 conference, Bergen, Norway (1979) and N. Schmitz, Proc. Int. Symp. on Lepton and Photon Interactions at High Energies (FNAL, Batavia, 1979).
81) J. Badier et al., Paper presented to the 1979 EPS Conference, Geneva, CERN EP preprint 79-67 (1979).
E. Berger, Invited talk at Orbis Scientiae 1979, Coral Gables, SLAC preprint PUB-2314 (1979).
Invited talks by W. Kienzle and J. Pilcher, Proc. Int. Symp. on Lepton and Photon Interactions at High Energies (FNAL, Batavia, 1979).
82) G. Parisi and R. Petronzio, Nucl. Phys. B154, 427 (1979). These results supersede those of Dokshitzer, D'yakonov and Troyan (Ref. 14) on the so-called T form factor, and are supported by the analyses of
C.L. Basham, L.S. Brown, S.D. Ellis and S.T. Love, Phys. Lett. 85B, 297 (1979).
C.Y. Lo and J.D. Sullivan, Univ. Illinois preprint ILL-(TH)-79-22 (1979).
83) See the second paper by Altarelli, R.K. Ellis and Martinelli in Ref. 73.
84) A.P. Contogouris and J. Kripfganz, Some quantum chromodynamic corrections to the Drell-Yan formula, McGill Univ. preprint (1979).
85) E. Berger and S.J. Brodsky, Phys. Rev. Lett. 42, 440 (1979). This possibly important source of scaling violation, together with uncertainties in the absolute normalization of the Drell-Yan cross-section and the importance of subasymptotic corrections, raise questions about the interesting scaling deviation analysis of
D. McCal et al., Phys. Lett. 85B, 432 (1979).
86) K.J. Anderson et al., Enrico Fermi Institute preprint EFI 79-34 (1979).
87) TASSO Collaboration, R. Brandelik et al., Phys. Lett. 86B, 243 (1979).
MARK J Collaboration, D.P. Barber et al., Phys. Rev. Lett. 43, 830 (1979).
PLUTO Collaboration, Ch. Berger et al., DESY preprint 79/57 (1979).
88) See the invited talks by
H. Newman, C. Berger, G. Wolf and S. Orito, Proc. Int. Symp. on Lepton and Photon Interactions at High Energies (FNAL, Batavia, 1979).
89) K. Koller and H. Krasemann, DESY preprint 79/52 (1979).
J.D. Jackson, preprint TH.2730-CERN (1979) argues for vector gluons on the basis of the ratio of hadronic decay widths of charmonium χ states. This argument seems precarious on account of the large radiative corrections apparently necessary in the QCD calculations.
90) J. Ellis and I. Karliner, Nucl. Phys. B148, 141 (1979).
91) Koller, Walsh and Zerwas, Ref. 71 and J. Ellis and F. Martin, unpublished.
92) R.K. Ellis and R. Petronzio, Phys. Lett. 80B, 249 (1979); see also
K. Konishi, A. Ukawa and G. Veneziano, Phys. Lett. 80B, 259 (1979). For other considerations on the perturbative widths of QCD jets in the framework of the Sterman-Weinberg analysis[19], see
I.I.Y. Bigi and T.F. Walsh, Phys. Lett. 82B, 267 (1979).
I.I.Y. Bigi, Phys. Lett. 86B, 57 (1979).
P. Binétruy and G. Girardi, Phys. Lett. 83B, 339 (1979).
G. Curci, M. Greco and Y. Srivastava, preprint TH.2632-CERN (1979).
M.B. Einhorn and B.G. Weeks, Nucl. Phys. B146, 445 (1978).
E.G. Floratos, Phys. Lett. 86B, 196 (1979).
W. Furmański, Ref. 67.
T. Gottschalk, G. Monsay and D. Sivers, Argonne preprint ANL-HEP-PR-79-16 (1979).
K. Shizuya and S.-H.H. Tye, Phys. Rev. Lett. 41, 787 (1978).
P.M. Stevenson, Phys. Rev. Lett. 41, 787 (1978); Nucl. Phys. B156, 43 (1979).
B.G. Weeks, Phys. Lett. 81B, 377 (1979). For a discussion of the oblateness of QCD jets, see
S.J. Brodsky, T.A. DeGrand and R.F. Schwitters, Phys. Lett. 79B, 255 (1978).
93) Present jets probably have non-perturbative widths, and it is unclear that the perturbative widths will be very visible soon.
W. Furmański, preprint TH.2664-CERN (1979).

94) Einhorn and Weeks, and Shizuya and Tye, Ref. 92.
95) C. Michael and L. van Ryckeghem, unpublished (1978).
96) C.L. Basham, L.S. Brown, S.D. Ellis and S.T. Love, Phys. Rev. D 17, 2298 (1978); Phys. Rev. Lett. 41, 1585 (1978); Phys. Rev. D 19, 2018 (1979); and paper cited in Ref. 82. See also
S.-Y. Pi, R.L. Jaffe and F.E. Low, Phys. Rev. Lett. 41, 142 (1978).
G.C. Fox and S. Wolfram, Phys. Rev. Lett. 41, 1581 (1978); Nucl. Phys. B149, 413 (1979); Phys. Lett. 82B, 134 (1979); CalTech preprint CALT-68-723 (1979). Early examples of QCD-motivated measures of hadronic final states which are not wedded to the jet structure are thrust and spherocity:
E. Farhi, Phys. Rev. Lett. 35, 1609 (1977).
H. Georgi and M. Machacek, Phys. Rev. Lett. 35, 1237 (1977).
97) R.J. DeWitt, L.M. Jones, J.D. Sullivan, D.E. Willen and H.W. Wyld, Phys. Rev. D 19, 2046 (1979).
98) W.A. Bardeen and A.J. Buras, Phys. Rev. D 20, 166 (1979).
99) Llewellyn Smith, also Brodsky, DeGrand, Gunion and Weis, Ref. 20.
100) For Drell-Yan production by photon beams, see:
L.M. Jones, J.D. Sullivan, D.E. Willen and H.W. Wyld, Univ. Illinois preprint ILL-(TH)-79-12 (1979).
For deep inelastic Compton scattering, see:
Tu T.-s. and Wu C.-m., preprint TH.2646-CERN (1979).
Chang C.-h., Tu T.-s. and Wu C.-m., preprint TH. 2675-CERN (1979).
For the QCD analogue process[77], see:
Tu T.-s., preprint TH.2690-CERN (1979).
For a proposed test of the three-gluon coupling by measuring the polarization of large p_T single photons, see:
A. Devoto, J. Pumpkin, W. Repko and G.L. Kane, Michigan State Univ. preprint MSU/HEP (1979).
101) G. Parisi, Phys. Lett. 84B, 225 (1979).
102) Brodsky and Lepage, Ref. 21. The first of these papers includes a discussion of the inclusive-exclusive connection which is to be extended and corrected in S.J. Brodsky and G.P. Lepage, to be published (S.J. Brodsky, private communication).
103) Brodsky, Frishman, Lepage and Sachrajda, Ref. 21. See also the remarks by Efremov and Radyushkin, Ref. 21, and Parisi, Ref. 101.
104) A. Duncan and A.H. Mueller, Columbia University preprint CU-TP-162 (1979) and discussion after this talk. Another discussion of baryon form-factors is given by
I.G. Aznauryan, S.V. Esaybegyan, K.Z. Hatsagortsyan and N.L. Ter-Isaakyan, Yerevan Physics Institute preprint 342(67)-78 (1979). The points raised by both these sets of authors will be discussed (and answered?) in Brodsky and Lepage, Ref. 102.
105) S.J. Brodsky and G.R. Farrar, Phys. Rev. Lett. 31, 1153 (1973); Phys. Rev. D 11, 1309 (1975).
V.A. Matveev, R.M. Muradyan and A.V. Tavkhelidze, Lett. Nuovo Cimento 7, 719 (1973). In QCD multiple scattering diagrams [P.V. Landshoff, Phys. Rev. D 10, 1024 (1974)] which threaten to violate dimensional counting rules in fact fall faster than any power of Q^2. See Brodsky and Lepage, Ref. 21.
106) P.G. Crabb, et al., Phys. Rev. Lett. 41, 1257 (1978).
107) G. Preparata and J. Soffer, preprint TH.2712-CERN (1979).
For a bizarre QCD-motivated mechanism to get large polarization, see:
S.J. Brodsky, C.E. Carlson, and H. Lipkin, SLAC preprint PUB-2305 (1979).
108) W. Furmański, R. Petronzio and S. Pokorski, Nucl. Phys. B155, 253 (1979).
A. Bassetto, M. Ciafaloni and G. Marchesini, Phys. Lett. 83B, 207 (1979).
K. Konishi, Rutherford Laboratory preprint RL-79-035 T.241 (1979).
109) D. Antreasyan and R. Petronzio, private communication (1979).
110) D. Amati and G. Veneziano, Phys. Lett. 83B, 87 (1979).
111) E. Reya, Phys. Rev. Lett. 43, 8 (1979).
112) I.I.Y. Bigi, preprints TH.2681-CERN and TH.2707-CERN (1979).
113) A. De Rújula, B. Lautrup and R. Petronzio, Nucl. Phys. B146, 50 (1978).
For the related idea of an asymmetry in heavy quark production at large p_T, see
R.W. Brown, D. Sahdev and K.O. Mikaelian, Ohio State Univ. preprint 88 (1979).
K.O. Mikaelian, R.W. Brown and D. Sahdev, Ohio State Univ. preprint 92 (1979).
114) Chicago Tribune, issue of Aug. 28, 1979, p. 1.

Speaker - A.H. Mueller - Columbia University

Q. Tony Duncan and I have re-examined the Brodsky-Lepage calculations and we find that for the pion form factor the subleading logarithms will organize themselves according to a renormalization group equation. However, for the nucleon form factor, by explicit example, we find logarithms at the non-leading level which cannot in fact be part of our renormalization group equation. So our conclusion is that the Brodsky-Lepage calculation is an asymptotic limit in an asymptotically free theory for the pion, but not for the nucleon. In general, the form factor is very complicated. For example, in ϕ^3 in six dimensions, no form factors will come from a renormalization group. For the pion, vector form factors would, but scalar form factors would not. So I think it is a very complicated situation.

A. So, if I understand you correctly, you say that there are apparently non-leading logarithms which actually dominate the leading logs?

Q. They are non-leading logarithms. I said they do not form part of a renormalization group equation. Whether they dominate or not is not known because one does not know how to group them yet.

A. I see, so you just say that the prediction of Brodsky and Lepage is uncertain for the nucleon form factor.

Q. At the moment it is a leading logarithm calculation for the nucleon. For the pion it actually is a limit of a renormalization group equation which holds to all logs.

Speaker - S.G. Matinyan - Yerevan Physics Institute

Q. I would like to make a remark on the theoretical calculations of elastic form factors. There are two calculations of nucleon form factors. One is due to Brodsky and Lepage, and the second is a calculation done in Yerevan, and they disagree. Unfortunately the results of Yerevan have not reached to the rapporteur but are circulating at the conference. The result is that unfortunately Brodsky and Lepage made a mistake where they relied on Leningrad results. They did not check the results when they separated the hard amplitude from the soft wave-function amplitude. They corrected some mistakes but also made an additional mistake. The result obtained by the Yerevan group is very strange, but if you take the non-relativistic wave function of the nucleon to be a point at the starting point of the approximation, you will obtain a negative sign for the proton form factor. It is very strange and indicates that we absolutely do not understand the non-relativistic nucleon structure. I would like you to pay attention to this fact. It is very important.

A. Okay. Thank you very much. If they send me that paper then I will report on it. I think that one point I should emphasize is that this work on the form factors does depend on assumptions about the hadronic wave function. The softness of it has not been justified in perturbation theory but depends on assumptions which go beyond it. I should emphasize that one of the reasons why I need speed is to escape from angry protesters about people I did not refer to.

Speaker - G. Wolf - DESY

Q. I have a question on gluestrahlung in e^+e^- annihilation. The basic diagram gives two coloured quarks plus a coloured gluon. In order to get singlet states you probably have to sum various diagrams. The question is, do you expect interferences after summation?

A. I do not quite understand your question. Do you mean if you go to the central region where the three jets join together? Is that your idea? Or do you mean actually out in the jets themselves?

Q. I mean in the jets themselves.

A. I think that at least in perturbation theory it looks like the jets fragment independently. There should not be any interference between them. However, the central region where the three of them are coupled together I think is a very unknown region and it would be very amusing to study it experimentally.

Q. Would that be in the Mercedes region?

A. That is right. That is the centre of the Mercedes star. That, I think, would be a very interesting region. There are other cases, for example at the ISR, where we have seen large p_T jets coming out. We think we understand maybe what happens out in these jets, but we do not really understand how they couple together. Now at the ISR it is very complicated, because you have got four jets all coming together. In e^+e^- when you have three, it is a simpler situation and perhaps easier to study.

Speaker - J. Thaler - Illinois

Q. There is a big correction to the cross-section in the Drell-Yan process. Why is there not a similar big correction in the deep inelastic scattering?

A. What is actually computed when you do these large corrections is the difference between, if you like, the renormalizations of the deep inelastic cross-section and of the Drell-Yan cross-section. Now whether there is a large renormalization in any individual one or in any individual other is a matter of definition. Different renormalization prescriptions give you different renormalizations here and there. What is invariant is the difference, and that is the thing which I was quoting.

Speaker - V. Telegdi - ETH Zurich (and Chicago)

Q. I am confused by the facts. We have heard both from you and I think from other people about this mystic factor kappa of order two that destroys the agreement between primitive Drell-Yan and experiment, taking experiment to be μ-pair production by pions. But not so long ago the celebrated Columbia-Fermilab and so-forth experiment was taken as a proof that colour exists because the requisite factor was present. Furthermore, the Chicago-Princeton experiment also seemed to agree with naive predictions. So, are we making progress or are we going backwards?

A. Well I think we are probably progressing side ways. If you look in the Oxford English Dictionary you will find the definition of a parton -- it is a Scottish dialect word for a "crab". And the quotation that they gave is "moving as a parton is wont to do, sideways". But I have a serious answer to this question if you will permit me. This is, as I understand it, the latest version of the comparison between the CDHS sea and the CFS cross-section. You will have to ask our distinguished summer-up why it is that now they disagree whereas before they agreed. As far as the CIT data is concerned, I understand that originally when they reported approximate consistency, that was when they were using an $A^{1.12}$ extrapolation (correct me if I am wrong). Anyway, if they use an A^1 correction, which is what one would expect from Drell-Yan, then they find this discrepancy. Now presumably if you are analysing in the context of the Drell-Yan model you should take A^1 rather than $A^{1.12}$. I should have perhaps have said that I think the clearest evidence to date probably comes from the NA3 Collaboration who did not know that they should find two when they found two, and then they were shocked to find they should have found it, if you see what I mean.