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Coherence Effects in Deep Inelastic Scattering

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Abstract:
We present a framework for deep inelastic scattering, with bound state properties in accordance with a QCD force field acting like a vortex line in a colour superconducting vacuum, which implies some simple coherence effects. Within this scheme one may describe the results of present energies very well, but one obtains an appreciable depletion of gluon radiation in the HERA energy regime.

I. Introduction

In this paper we will consider the emission of gluonic bremsstrahlung in connection with deep inelastic scattering (DIS). We will be particularly interested in situations where there are possible observable effects from coherent emission as well as bound state properties.

In the original parton model the field quanta of the initial hadron are treated as independent quantities. Thus all dynamical results may be obtained by adding the contributions. The parton densities then occur as multiplicative flow factors. This picture corresponding to the wellknown Weizsäcker-Williams approximation in QED is particularly noticeable in the stochastical process schemes for parton branching, which are built on the Altarelli-Parisi equations.

This basically incoherent scheme is probably often a good approximation. Nevertheless it is possible to disentangle regions in phase space where coherence effects in the emission are important. One such region is the long wavelength limit. In that case coherence is at least partially taken into account by the prescription of "angular ordering" [1]. It is wellknown that Monte Carlo models with this feature reproduce experimental results better than those without it [2]. (This is primarily the case for cluster fragmentation models. With string fragmentation the difference is very small as the soft gluons here give very few extra hadrons.) Besides the prescription of timelike cascades ("final state radiation") there are corresponding prescriptions of coherence character for the spacelike cascades ("initial state radiation") [2,3].

It has also been pointed out that there should be a depletion of gluons with small x_B values due to the spacetime overlap of the long wavelengths [4].

In the opposite extreme it is possible to imagine that in the emission of very hard gluons close to the kinematic boundary, there may, depending upon the structure of the system, be modifications of the allowed phase space regions due to coherence effects.

We will in this paper consider the consequences for the gluon radiation in case the hadronic state and more generally all confined colour force fields behave like vortex lines in a colour superconducting vacuum. Such a system has basically a onedimensional ordered extension with similarities to a linked dipole chain (fig. 1). Although the vortex core is surrounded by a field with an expected (transverse) extension of about a fm, the main energy is concentrated to the core. This implies that the dynamics of the vortex line is that of a massless relativistic string with the valence charges (q-, $\bar{q}$- or (effective) qq-particles) at the endpoints.

We note that an undisturbed colour singlet system in its ground state does not emit any radiation. If, on the other hand, one of the charges at the endpoint is accelerated by an outside, e.g. electromagnetic, field pulse, then there will be colour separation and subsequent radiation from the gauge field.

It is instructive to consider a bound system in electrodynamics, like e.g. an e^+e^- system, when one of the charges, e.g. the e^-, is in a similar way accelerated. We will first neglect the internal motion of the pair and the binding force. Then the classical motion is as indicated in fig. 2a. The radiation can be described in two equivalent ways. One way is to say that the bound system is neutral and therefore does not radiate before $t=t_0$. For times t larger than t_0 the e^+ and the e^- move apart and then they will emit dipole radiation. An alternative way is to say that the e^+ is never accelerated, and therefore does not radiate, while the e^-, which changes direction, is the one to emit radiation.

The two descriptions correspond to the Feynman diagrams in figs. 3a and 3b respectively. In both cases the amplitude has two contributions which in general will interfere. The separation is however gauge dependent and it is e.g. possible to choose a gauge such that only the upper two diagrams give contributions. Then in both cases the photon emission seems to come from e^- "after" it has been struck. Thus in this particular gauge the two descriptions become equal.

In case the bound state properties are not negligible, then the classical trajectories will be as in fig. 2b. We note that although the e^+ and e^- are accelerated all the time in the bound state there is quantum mechanically no radiation as long as the system is in its ground state, i.e. "before" the interaction. Then it is likely that the first alternative, where the radiation is determined by the charge separation after the interaction, should provide a more accurate description. For the second alternative it will be necessary to include in the amplitude a contribution corresponding to the transition from the bound state wave function to a precise value of the electron momentum as well as a contribution from the e^+ which is also accelerated in this case. All these contributions may interfere and should be added coherently. The contribution from the localization of the e^- in momentum space has similarities to photon emission when an electron is captured in inverse β decay. That process corresponds to a similar localization in coordinate space.

For the corresponding situation in QCD we expect several differences. Firstly the force field itself contains charge, [illegible] charges [illegible] accelerated (at least after some time) as in fig. 2b. In our ordered dipole chain-string field this acceleration occurs in a coherent fashion. Secondly, the force field of the bound state has its total energy distributed over a (for the vortex line case one-dimensional) region in space. Any emission will involve only a part of the system (an ordinary radiation condition means a region of the order of a fraction of the wavelength) and consequently only that part of the energy will be available for the emission. We discuss this feature in detail in section 2 for the vortex line string hadron.

We will in this paper make use of the dipole approximation for the gluonic radiation. Our group has in a set of papers [5] investigated a somewhat different approach to the bremsstrahlung, viz. to take the dipole emission as such as the basic feature. We note that a basic difference between QED and QCD is that the emitted gluonic radiation in QCD carries charge by itself, thereby changing the originally radiating system. In case a quark and an antiquark go apart and emit a gluon, then the original "dipole antenna" is changed into two dipoles [6,5], which each may emit inside the corresponding characteristic cones. This coherence condition which is built into the dipole treatment has been incorporated in the different other schemes by means of the above mentioned angular ordering prescriptions.

One feature of *DIS* which it is necessary to treat in detail is the state properties in case a nonvalence quark or antiquark is struck out. This has in e.g. the earlier Lund group treatments been handled in a rather cavalier fashion by some simple parametrizations [7]. In the high energy Hera region the sea-quark interactions constitute the major part of the cross-section and we feel that it may be worthwhile to treat these interactions in a more precise way according to string dynamics. The ensuing model based upon the unique breakup probability of a Lund string is presented in section 3.

Section 4 contains a discussion of the parameters of the model determined from the present energy regime as well as the prediction for the Hera energy regime and section 5 contains some concluding remarks on the soft coherence scheme presented in this paper.

II. Bremsstrahlung from an Extended Colour Source

In this section we will develop a model for the soft coherent radiation that we expect is emitted from an extended colour source of a vortex line character. We start with a brief review of the properties of the dipole approximation in an e^+e^--annihilation reaction [5]. The colour separation when two colour connected partons with (massless) energy momentum vectors p_1 och p_2 go apart corresponds to a dipole antenna with total cms energy

$$\sqrt{s} = \sqrt{(p_1+p_2)^2} \tag{1}$$

This may emit gluonic bremsstrahlung transforming the system to three partons with (massless) energy momentum vectors k_j, j=1,2,3 according to the dipole formula:

$$d\sigma = \frac{3\alpha_s}{4\pi} \cdot \frac{x_1^2+x_3^2}{(1-x_1)(1-x_3)}\, dx_1 dx_3 \simeq \frac{\alpha_0}{\ln(k_T^2/\Lambda^2)} \frac{dk_T^2}{k_T^2}\, dy \tag{2}$$

with the variables

$$\begin{aligned} s_{12} &= (k_1+k_2)^2 = s(1-x_3) \\ s_{23} &= (k_2+k_3)^2 = s(1-x_1) \\ k_T^2 &= \frac{s_{12}s_{23}}{s} \\ y &= \frac{1}{2}\ln(s_{12}/s_{23}) \end{aligned} \tag{3}$$

The constant α_0 is defined by the relation

$$\frac{3\alpha_s}{2\pi} = \frac{6}{(11-2/3\cdot N_f)\ln(k_T^2/\Lambda^2)} = \frac{\alpha_0}{\ln(k_T^2/\Lambda^2)} \tag{4}$$

while a_1 and a_3 take on the values 2 and 3 depending upon if the parton j=1 or 3 is a (anti)quark or a gluon [8]. The variables k_T and y have in the long wavelength limit the meaning of transverse momentum and (pseudo-)rapidity of the emitted gluon 2 with respect to the original axis in the dipole cms. Then the kinematical range is:

$$|y| \leq \ln(\frac{\sqrt{s}}{k_T}) \tag{5}$$

which may be described as the interior of a triangle in the y-$\ln(k_T)$ plane.

The system (k_1,k_2,k_3) is oriented with respect to the original system (p_1,p_2) in such a way that the colour flow in and out of the system is minimally changed by the emission. Thus in the (p_1,p_2) cms system one chooses

$$k_{T1}^2 + k_{T3}^2 = \text{minimum} \tag{6}$$

if (1,3) both correspond to gluonic emittors, while in a (q,g) or ($\bar{q}$,g) emission system one chooses the quark or antiquark to take the recoil alone. The azimuthal angle around the dipole axis is chosen at random.

In this way one obtains a scheme which may be repeated again with the partons (k_1,k_2) and (k_2,k_3) forming new dipoles. Recent investigations [9] have shown however, that small modifications in the scheme close to the kinematical boundaries give noticeable differences at high energies (LEP, CLIC) even if the different schemes are tuned to reproduce experimental data in the PEP-PETRA range. Thus if the full phase space of eq. (5) is used in each dipole the total hadronic multiplicity is enhanced by ~ 25 % at LEP or SLC compared to a scheme in which the k_T's are ordered such that the k_T of one dipole is always smaller than the k_T of the previous one. This latter k_T-ordered scheme gives results very close to the other Monte Carlo schemes on the market, e.g. the Lund JETSET 6.3 [10] or the Webber Monte Carlo [11]. Using the full phase space means in some sense that each dipole decays (except for the recoils) independently of everything else. We note that perturbative QCD can not be used as a guide when choosing the best scheme, because in these regions close to the kinematical limits, the tree approximation, which implies a classical branching process, is not trustworthy. Nevertheless, although the two schemes give different results for very high energy e^+e^--annihilation, when we in the following apply these ideas to deep inelastic lepton scattering the differences are very small, both at EMC and at HERA energies.

We now move on to deep inelastic lepton scattering and consider the situation when a vortexline (dipole-chain, stringfield) hadron interacts with an external field source. In that way the endpoint of the string which is assumed to correspond to a flavour charged valence quark obtains a momentum transfer. The discussion of the events when sea quarks are involved will be postponed to the following section.

In the final state hadronic cms the interaction will transform the initial hadron into a state where the hit endpoint is moving in one direction with energy W/2 and an extended remaining system with a string moving in the opposite direction, also with energy W/2 (see fig. 4a). We note that this system is Lorentz contracted so that its extension in the longitudinal direction is small.

The ensuing colour separation means that we will emit radiation, e.g. a gluon with the transverse momentum k_T and the rapidity y. It is possible to make a longitudinal Lorentz transformation to a frame in which the gluon is emitted at 90^o (cf fig. 4b). We observe that the coherence condition for radiation from an extended source means that only a fraction of the source will in general be involved. Thus the emission from an antenna is reduced when the antenna size is larger than the wavelength. In particular if the wavelength of the emitted radiation is $\lambda=2\pi/k_T$ then only a region of the dipole chain string with transverse extension approximately equal to $\lambda/2$ will give constructive interference.

In this way we obtain a kind of "effective" antenna or effective dipole which consists of the struck endpoint and a part of the remainder system of the order of a fraction of the wavelength. If this part carries a fraction $a(k_T)$ of the full remainder energy then the momenta of this effective dipole are given by (in the Lorentz frame of fig. 4b)

$$\begin{cases} \text{endpoint:} & e^{-y}\cdot W/2 \\ \text{part of remainder:} & a(k_T)\cdot e^{y}\cdot W/2 \end{cases}$$

For such a dipole the kinematic limit is given by

$$k_T < \frac{Wa(k_T)}{e^{y}a(k_T)+e^{-y}} \qquad (7)$$

The spatial extension of the effective dipole should be approximately given by $\lambda/2=\pi/k_T$. Thus if the energy of the remainder system is distributed evenly along the string, then we expect the fraction

$$a(k_T) = \mu/k_T \qquad (8)$$

to be involved with μ a parameter related to the inverse of the hadronic size, i.e. μ should be of the order of a hadronic mass. One could also imagine that *the energy is more (or less) concentrated close to the endpoint*. We have consequently investigated the more general case

$$a(k_T) = (\mu/k_T)^{\alpha} \qquad (9)$$

for various values of the parameter α. In section 4 we find that a good description of EMC data is obtained with α=1, i.e. with the distribution in eq. (8).

If we insert the relation in eq. (8) or (9) into eq. (7) we see that the consequences of an extended colour charge distribution is that hard gluons in the target fragmentation region and in the central region are suppressed. For large values of y ($y>\alpha/2\ \ln(k_T/\mu)$) the limit in eq. (7) agrees approximately with the kinematical limit in eq. (5). These gluons get their energy mainly from the pointlike struck quark and they are therefore not sensitive to the colour charge distribution in the target remnant. For central or backward-moving gluons with $y < \alpha/2\ \ln(k_t/\mu)$ the limit is approximately given by the expression

$$\ln\left(\frac{k_T^{\alpha+1}}{W\mu^{\alpha}}\right) \leq y \qquad (10)$$

The allowed region is shown in fig. 5. In particular we see that there is a maximal k_T given by

$$k_{T\max} \sim (W^2\mu^{\alpha})^{\frac{1}{\alpha+2}} \qquad (11)$$

Naturally the expressions in eq. (8) or (9) are only applicable when $k_T>\mu$. For longer wavelengths, such that $k_T<\mu$ the whole remainder vortex line acts coherently. When $k_T>\mu$ we expect that the part with energy fraction $a(k_T)$ which is involved in the gluon emission also gets a recoil while the rest is unaffected by the emission. This recoil causes a kink on the string and thus acts as an extra gluon.

It is interesting to compare our suppression of gluons in the target fragmentation region with the corresponding suppression in the normal treatment of initial state radiation. In that approach a quark, which initially has a large momentum fraction x', can emit a gluon, thereby reducing its momentum fraction to x_B. The relative probability for this process is given by the ratio of the structure functions

$$\frac{f(x',Q_0^2)}{f(x_B,Q^2)} \sim \left[\frac{1-x'}{1-x_B}\right]^{\rho} \qquad (12)$$

Here ρ is a power of order 3-4. High energy gluons correspond to large values of x' which are strongly suppressed.

For a gluon with transverse momentum k_T and rapidity y (in the hadronic cms) it is easily seen that

$$\frac{1-x'}{1-x_B} = 1 - \frac{z_-}{1-z_+} \qquad (13)$$

where $z_{\pm} = \frac{k_T}{W} e^{\pm y}$

are the fractional lightcone momenta for the gluon.

A fixed value for the suppression factor in eq. (12) corresponds to the dashed curve in fig. 5. Apart from scaling violations in the structure functions this curve is scaling and thus $k_{Tmax} \sim W$. Thus although the two suppression mechanisms turn out to be similar at low energy, there is evidently an increasing difference at higher energies.

III. Sea Quark Interactions

In the previous section we have only discussed the situation when a valence quark has been struck by a virtual photon. It is necessary to study also sea quark interactions. As the energies increase, these interactions correspond to a larger and larger part of the cross section.

When the struck valence quark is replaced by a sea quark the remnant of the initial proton (in case of a lepton-proton scattering) is changed from a diquark system to a system containing the three original valence quarks plus a partner of the struck sea quark. This may affect both the gluon radiation and the fragmentation of the partonic state into hadrons.

In earlier versions of the Lund model [7] the usual parton model picture is adopted in which the struck quark radiates while the parton remnants are only spectators. Thus the only modification of the gluon radiation in case of a struck sea quark is caused by possible differences in the structure functions for the quarks which can start the initial state radiation cascade. The fragmentation is described in terms of two independent string pieces $[q_s (qq)_v]$ and $[\bar{q}_s q_v]$ with a rather ad hoc parametrization of the energy sharing between them.

Which modifications of the scheme described in section II should we expect when a struck valence quark is replaced by a sea quark? To be specific we consider a struck quark and not an antiquark. Here the gluons are emitted by a dipole spanned between the struck quark which is an essentially pointlike colour 3 charge, and the target remnants, for which the colour $\bar{3}$ charge is confined to a region of hadronic size. In case of a sea quark interaction we expect the initial $q_s\bar{q}_s$ pair to be confined within the same hadron size as the valence quarks. Thus the distribution of the colour $\bar{3}$ charge in the target remnant is very similar for the $q_v q_v q_v \bar{q}_s$-system in a sea quark interaction and the $q_v q_v$-system in a valence quark interaction. Therefore we expect that the gluon emission is essentially the same in the two cases.

When the struck quark moves away the vacuum compresses the colour electric field to a linear structure like a vortex line. When sufficient energy is stored in this field it can break via $q\bar{q}$-pair creation. The length of the field when it breaks is larger than a hadronic radius (in the rest frame of the produced hadron). Therefore this string ought to be rather unaffected by finer details of the charge distribution in the target remnants, and consequently also the momentum distribution of the final state hadrons ought to be similar.

It is wellknown that there is a unique breakup probability [12] for a Lund string, i.e. given the flavour f_0 at the endpoint and the (anti)flavour $\bar{f}_1$ at the breakup (determined by tunnelling arguments) a hadron $f_0\bar{f}_1$ (transverse mass m_T) will take a fractional energy momentum z in accordance with the formula

$$\frac{dP}{dz} = N \frac{dz}{z} (1-z)^a \exp-(bm_T^2/z) \qquad (14)$$

Here N is a normalization constant, b is a flavour independent parameter while the parameter a may be flavour dependent although in most applications it has been assumed to be independent of flavour. For a three quark state a qq-system acts as a colour $\bar{3}$ at one end of the string.

According to the discussion above we expect that also for sea quark interactions the breakup probability is governed by the fragmentation function of eq. (14) in particular also for the first rank hadron in the target fragmentation end. To give a precise structure to the first rank hadron we assume that it will carry the flavours $q_v \bar{q}_s$.

This is further supported by a study of the classical motion of the system. The $q_s \bar{q}_s$ pair is not likely to be in a colour singlet state. In an Altarelli-Parisi evolution picture they originate from a gluon which is a colour octet and in a string picture they correspond to a virtual break of the string. The struck sea quark q_s will then be linked to a diquark $(qq)_v$. The remaining system $q_v \bar{q}_s$ can not directly form a hadron because its invariant mass is too small. It must hang on to the $(qq)_v$ (which is pulled away by the q_s) until the confining force has sufficiently increased its momentum (as measured in the lab frame). Therefore the whole system resembles a single string as illustrated in fig. 6a.

The net result of our considerations is that the size of the colour charge distribution in the target remnant system and the confining force field are very similar in sea quark and valence quark interactions. Consequently also the radiation of gluons and the momentum distribution for the hadrons are approximately the same.

We have here discussed the case of a kicked out sea quark. In case of a sea antiquark the only modification is that the first rank hadron (in the target fragmentation region) becomes a baryon (cf fig. 6b).

A comparison of the hadronization in the target fragmentation region between this scheme and the earlier treatment [7], shows that the difference is very small as long as u, d and s flavours are involved. There is, however, a set of interesting possibilities in case charm or bottom quantum numbers can be excited. It takes quite some sophistication in the detector setup however to disentangle such a sigral and we will not pursue this question here.

IV. Comparison with Experimental Results and Predictions for Higher Energies

In this section we will compare our method to experimental data from EMC [14], and determine the parameters of the theory, as well as give predictions for larger energies, in particular for the HERA energy regime.

To simulate the gluon radiation in accordance with the discussion in section II we will make use of the dipole cascade Monte Carlo, Ariadne 2 [13]. This program describes very well the results of e^+e^--annihilation events in the PETRA-PEP energy regime. In addition to the two parameters α and μ in our coherence condition Ariadne in principle contains a single parameter Λ_{QCD}=250 MeV, but as in all the parton cascade descriptions there are a set of kinematical choices and conventions. Basically the probability for emitting a gluon (eq. (2) above) is used repeatedly together with a prescription for the calculation of the new dipoles discussed above.

To calculate the cross section as a function of Q^2 and x_B we use the computer program LEPTO [7] with structure functions given by EHLQ1 [15]. The fragmentation of the final state string system is done using the Lund fragmentation scheme (at present JETSET 6.3 [10] with the parameters chosen from the comparison to e^+e^--annihilation).

In order to determine the parameters α and μ we have used the EMC data [14], in particular distributions related to the transverse momentum activity. In figs. 7-9 we show comparisons with data for the seagull effect ($\langle p_T^2 \rangle$ as a function of x_F), the compensation of p_T for trigger particles with large x_F and the p_T-distributions in and out of the event plane. These distributions are most sensitive to α and μ but all other results are well described by our scheme.

It turns out that values of $\alpha \sim 0.75$ to 1.1 with values of μ between 0.6 and 1.0 GeV (there is a correlation so that small α corresponds to small μ for best results) are all compatible with the data using a "primordial p_T" $\langle p_{Tprim} \rangle \sim$ 0.0-0.4 GeV/c (cf table I for best values). We note that all these values have a very reasonable size ($\alpha=1$ corresponds to an even energy distribution along the string and $\mu \sim m_\rho$ to m_p corresponds reasonably to the hadronic mass scale).

When it comes to much larger energies in the Hera regime an inspection of fig. 5 immediately reveals that we expect a strong decrease in gluonic emission in the central region, in particular between the centre up to the fragmentation region when comparing with the usual approach. As mentioned above we have $k_{Tmax} \sim W^{2/3}$ (for $\alpha=1$) instead of $k_{Tmax} \sim W$. In the backward region there is instead a (small) increase.

This also comes out from a detailed MC study. In figs. 10-11 we exhibit the inclusive charged particle rapidity distribution as well as dp_T/dy from our model (full line) and the wellknown LEPTO MC [7] (dotted line) for Hera energies. The difference between the models is largest in the region $0 < y < 2$, and in fig. 12 we show the E_T-distribution for this region. As we expect from the rapidity distribution there is considerably more transverse energy flow in the region according to the LEPTO model. This is presumably one of the most direct ways to experimentallly distinguish between the two models.

The comparisons are done for $\alpha=1.0$, $\mu=m_p$, $\langle p_{Tprim} \rangle = 0.3$, but we have tested that the differences are basically the same for values of α, μ and $\langle p_{Tprim} \rangle$ inside our "allowed region" (table I).

It is evident that our coherence condition, which using the variable z_- from eq. (13) is

$$a(k_T) - z_- > 0 \qquad (15)$$

means a rather drastic "complete cutoff" for hard gluonic radiation. If we instead of the step function cutoff in eq. (15) introduce a smoother suppression factor like

$$\frac{1}{1+\left[\frac{z}{a(k_T)}\right]^{\beta}} \qquad (16)$$

we find that the difference between the models persists, even for very small values of β (the dashed line in figs. 10-12 are the results for $\beta=1$).

V. Concluding Remarks

In this paper we have presented a possible scenario where the bound state properties and a coherence condition together lead to a large depletion of gluon emission in the HERA energy range and above. We note that our model is built to provide a good description of the results from the presently available energies, and that all our results are based upon radiation and fragmentation properties in accordance with other models in e^+e^--annihilation. One may nevertheless question the stability of such model results. We note that the effects are basically the same for all α, μ and primordial p_T, which are compatible with the EMC-results, and they are also independent of whether the coherence condition is implemented by a sharp or a smooth cutoff.

We finally note that our approach has the particular advantage that the full coherence of the basic gluon radiation matrix element is incorporated. Therefore there is a smooth and simple transition between initial and final state radiation. This is difficult to incorporate in the ordinary approach where each parton radiates by itself [16].

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Table I

The values of α, μ and $\langle p_{Tprim}\rangle$ that give good agreement with the EMC-data [14] are strongly correlated. The parameter α may vary between 0.75 and 1.1 and the corresponding best values of μ and $\sigma=\langle p_{Tprim}\rangle$ are given in this table. For the results presented in figs. 7-9 we have used $\alpha=1$, $\mu=m_p$ and $\sigma=\langle p_{Tprim}\rangle=0.3$ GeV.

α	μ (GeV)	σ (GeV)
1.1	1.0	0.3
0.75	0.6	0.4

Figure Captions

1 The field of a vortex line in a type II superconductor is the same as the field from a chain of dipole links.

2a Classical motion of an e^+e^- system when bound state effects are neglected. At time $t=t_0$ the electron is accelerated by an external force. The bremsstrahlung may be described either by the e^+ and e^- going apart after $t=t_0$ or by the acceleration of the e^-.

b Classical motion of a bound e^+e^- system. After the interaction both the electron and the positron are accelerated. The bremsstrahlung is in this case more easily described by the charge separation after the interaction.

3a Feynman diagrams for the emission of a photon from a separating e^+e^- pair.

b Feynman diagrams for the emission of a photon from an accelerated electron.

4a A colour dipole is formed between a quark, which is hit by a virtual photon and the stringlike hadron remnants. This dipole can emit gluon radiation.

b The same process in a frame where the gluon is emitted at 90^o. Due to *destructive interference, only* a fraction of the hadron is involved in the emission.

5 The kinematically allowed region in the y-$\ln(k_T)$ plane. The dashed line is the upper limit for gluon radiation in our model (eq. (10)). The dotted line is a line of equal suppression in the ordinary parton model. At lower energies (fig. a) the difference between the two approaches is less important, but at higher energies (at HERA, fig. b) our treatment gives a strong suppression of hard gluon emission *compared* to the ordinary parton model.

6 The string that is formed when a sea quark (a) or an sea antiquark (b) is kicked out by a photon.

7 $<p_T^2>$ as a function of x_F in the hadronic cm frame. The full line is the result for $\alpha=1$, $\mu=m_p$ and $<p_{Tprim}>=0.3$ GeV. For comparison we show the result for $\mu=0.5$ GeV with the same values of α and $<p_{Tprim}>$ (dotted line). Data are from EMC [14]

8 Transverse momentum balance dp_T^{bal}/dy for trigger particles with $x_F>0.5$ (a) and with $-0.5<x_F<-0.2$ (b). Notation as in fig. 7.

9 Distribution of Σp_T^2 out of the event plane (a) and in the event plane (b). Notation as in fig. 7.

10 Charged particle rapidity distribution in the hadronic cm frame for neutral current events at HERA. $\alpha=1$, $\mu=m_p$ and $<p_{Tprim}>=0.3$ GeV. The full line corresponds to a "complete cutoff", i.e. eq. (15) or $\beta \to \infty$ in eq. (16), and the dashed line is for $\beta=1$. The dotted line is the result for LEPTO [7].

11 p_T-weighted rapidity distribution dp_T/dy for charged particles at HERA. Notation as in fig. 10.

12 Distribution of ΣE_T at HERA for charged particles with rapidity $0<y<1$ (a) and $1<y<2$ (b). Notation as in fig. 10.

FIG. 1

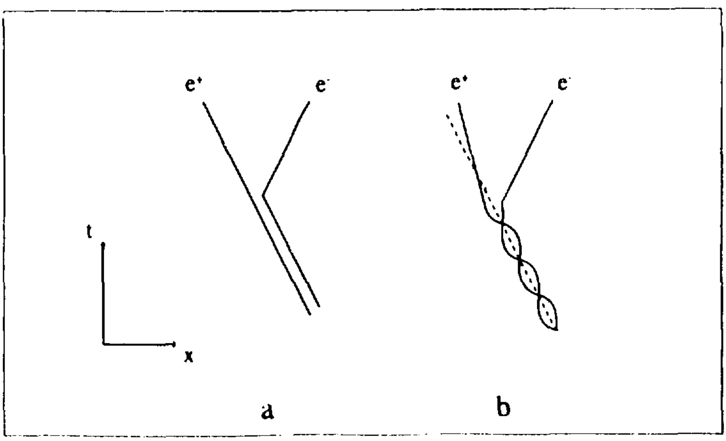


fig. 2

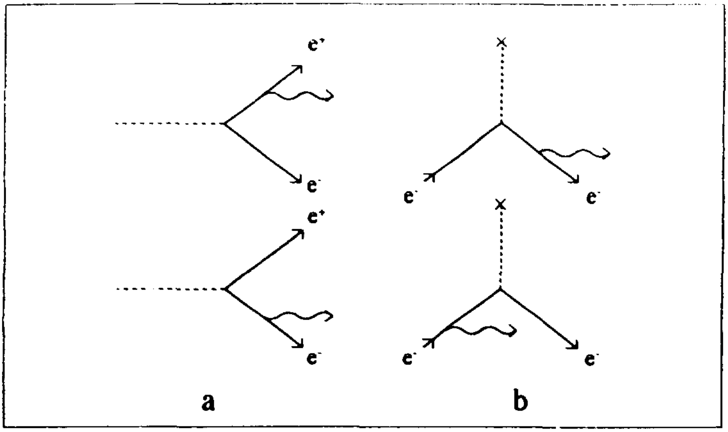


fig. 3

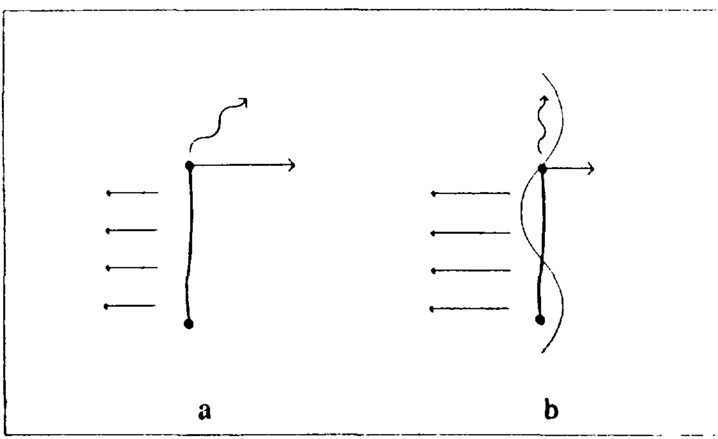


fig. 4

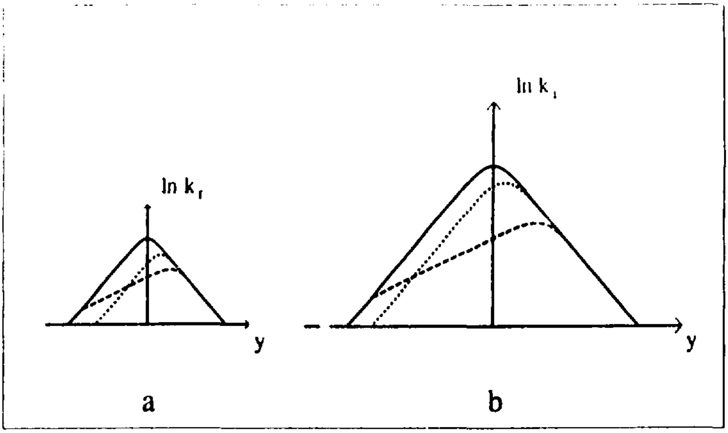


fig. 5

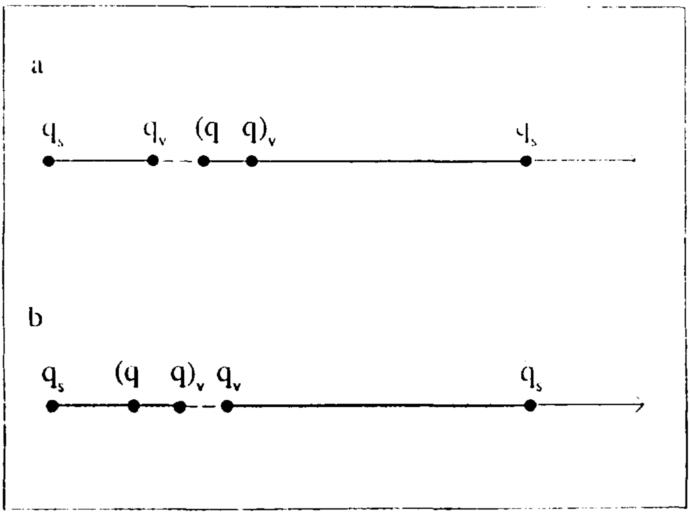
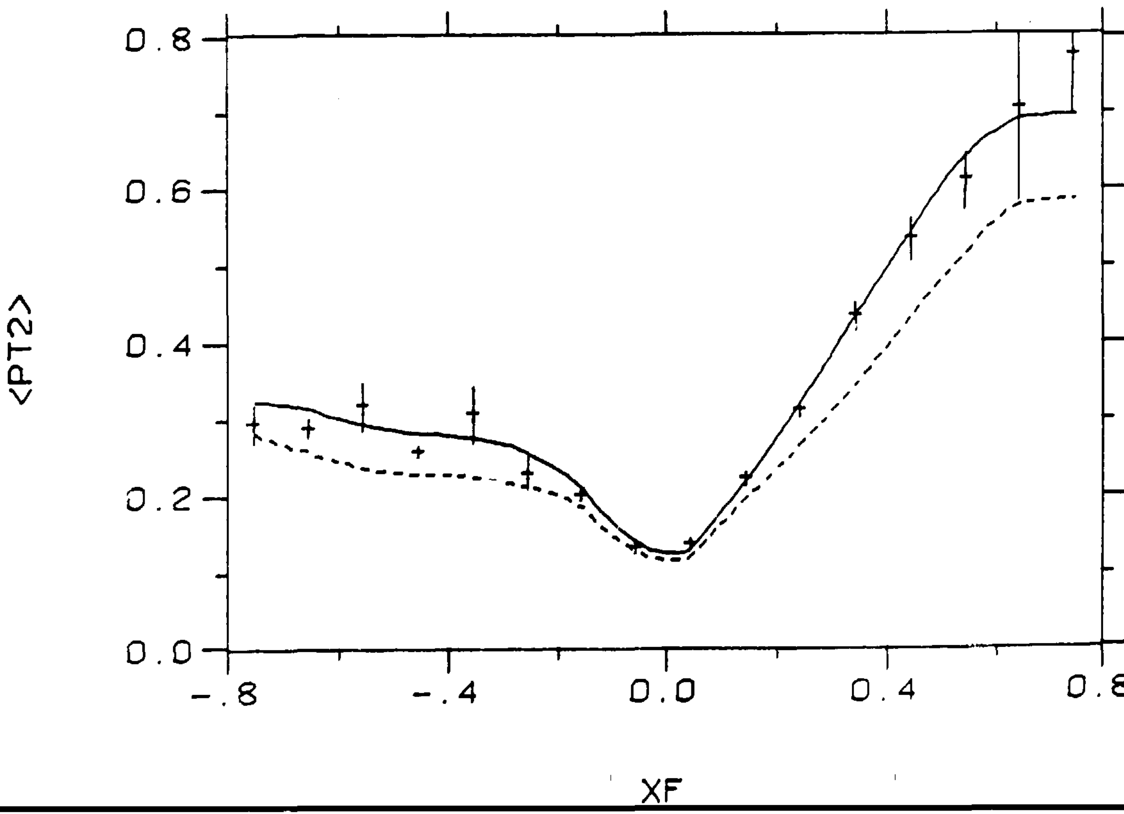


fig. 6



<PT2>
0.0
0.2
0.4
0.6
0.8
-.8
-.4
0.0
0.4
0.8
XF

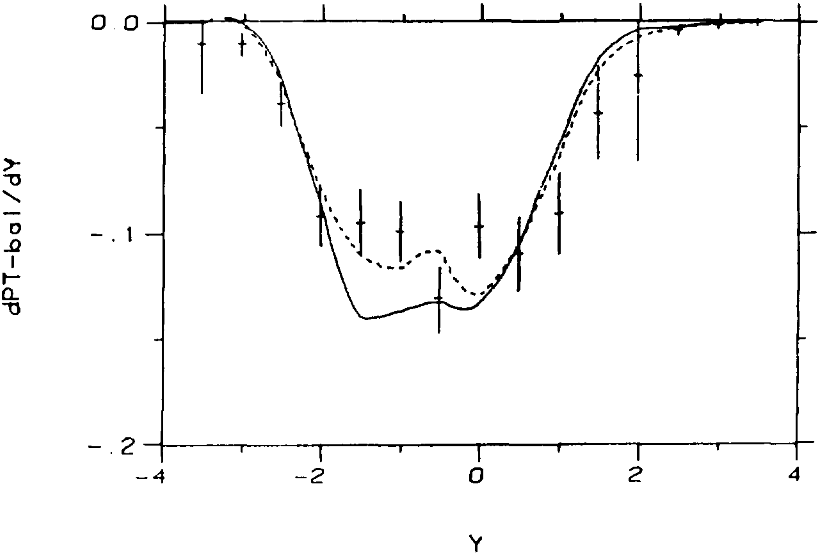


Fig 8a

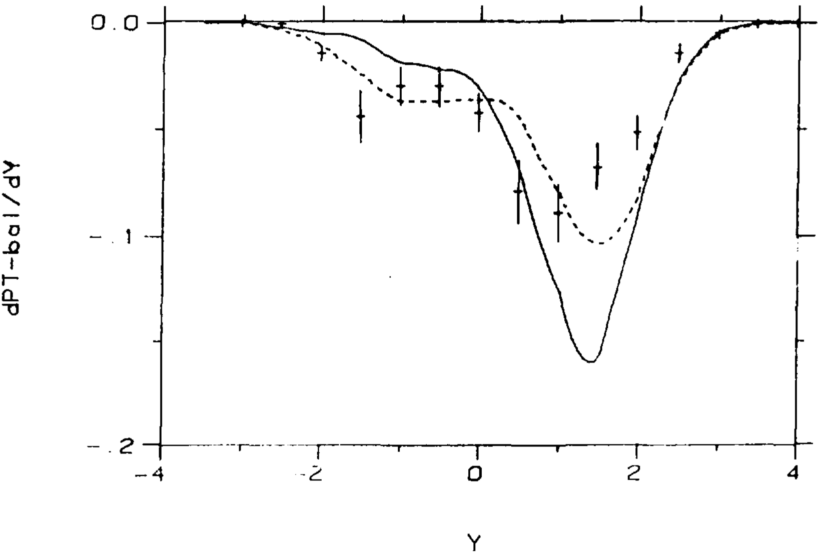
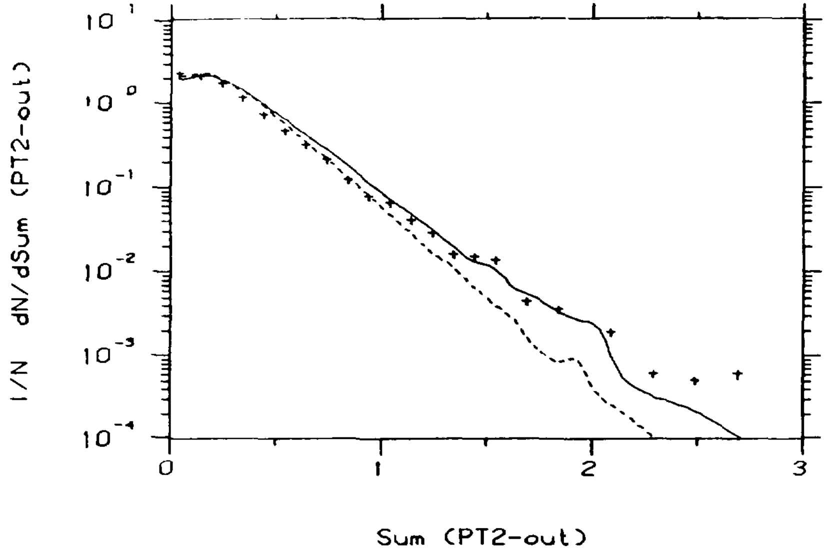
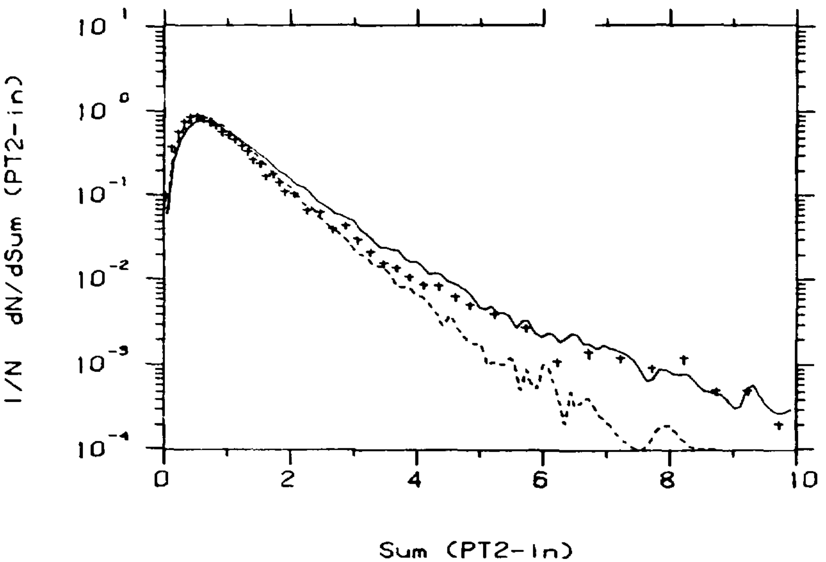


Fig 8b


10 1
10 0
10 -1
10 -2
10 -3
10 -4
1/N dN/dSum (PT2-out)
0
1
2
3
Sum (PT2-out)
Fig 9a


10 1
10 0
10 -1
10 -2
10 -3
10 -4
1/N dN/dSum (PT2-in)
0
2
4
6
8
10
Sum (PT2-in)
Fig 9b

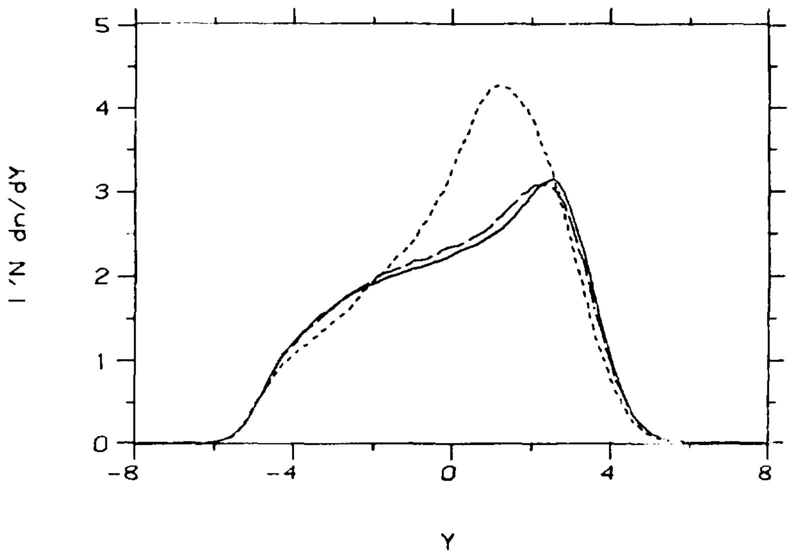


Fig 10

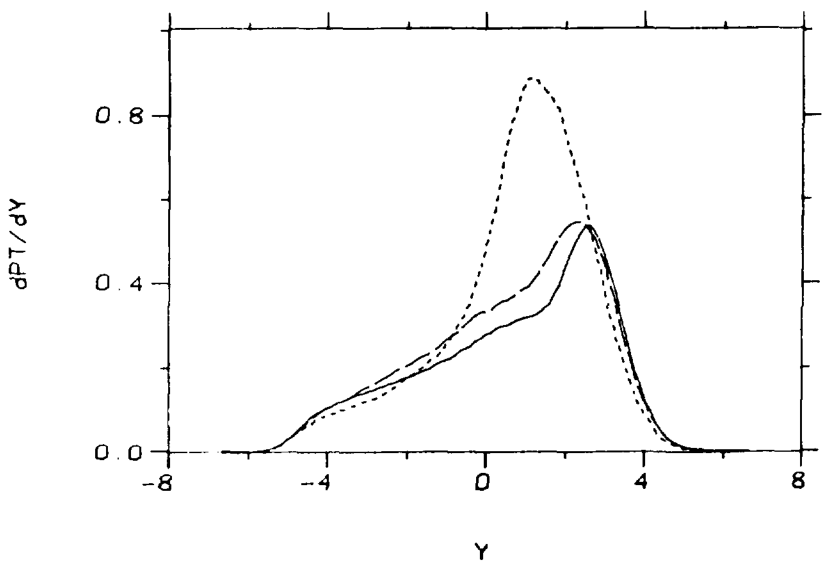
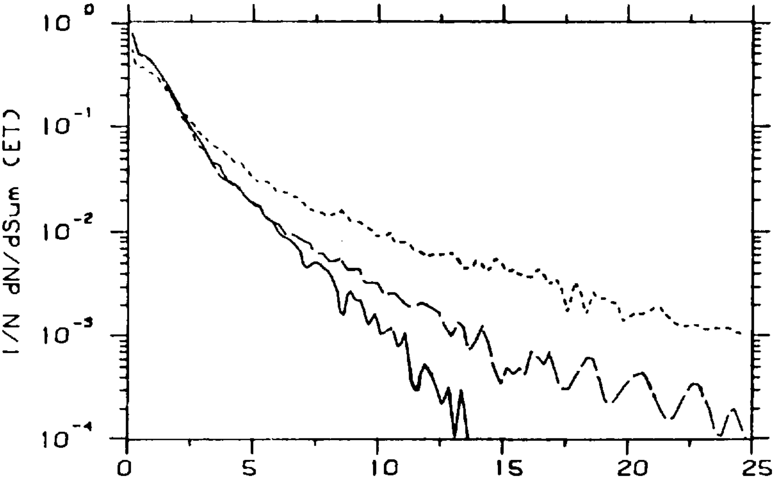
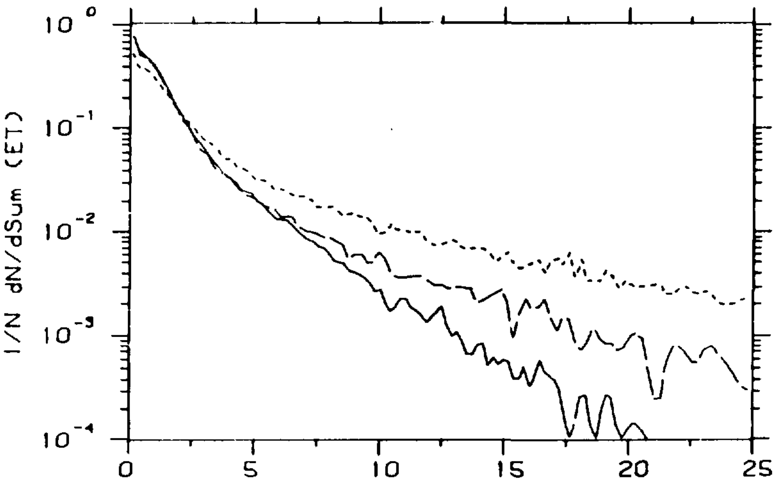


Fig 11



Sum (ET) for 0 < Y < 1

Fig 12a



Sum (ET) for 1 < Y < 2

Fig 12b