The reason why some data on the production of $\psi$- and $\Upsilon$-states disagree with QCD predictions, occasionally by well over one order of magnitude, is that the traditional method for performing the perturbative calculation of the cross section is simply wrong. The key mistake is to require that the heavy quark pair forms a color singlet at short distances, given that there is an infinite time for soft gluons to readjust the color of the $c\bar{c}$ pair before it appears as an asymptotic $\psi$ or, alternatively, $D\bar{D}$ state. We suspect that the same mistake is made in the description of rapidity gaps, i.e. the production of a color-neutral quark-antiquark pair, in terms of the exchange of a color neutral gluon pair. The $\psi$ is after all a color neutral $c\bar{c}$ pair and we will show that it is produced by the same dynamics as $D\bar{D}$ pairs; its color happens to be bleached by soft final-state interactions. This approach to color is suggestive of the unorthodox prescription for the production of rapidity gaps in deep inelastic scattering, proposed by Buchmüller and Hebecker. When applied to the formation of gaps between a pair of high transverse momentum jets in hadron collisions, the soft color approach suggests a formation rate of gaps in gluon-gluon subprocesses which is similar or smaller than in quark-quark induced events. Formation of gaps should increase when increasing transverse momentum or lowering energy, in contrast with 2-gluon exchange Pomeron models.

1 Introduction

The conventional treatment of color, i.e., the color singlet model, has run into serious problems describing the data on the production of charmonium and upsilon states. Specific proposals to solve the charmonium problem agree on the basic solution: its production is a two-step process where a heavy quark pair is produced first. At this stage perturbative diagrams are included whether the $c\bar{c}$ pair is color singlet or not. This is a departure of the textbook approach where only diagrams where the charm pair is in a color singlet are selected. In the Bodwin-Braaten-Lepage (BBL) formalism, the subsequent evolution of the pair into a colorless bound state is described by an expansion in powers of the relative velocity of the heavy quarks in theonium system. An alternative approach, color evaporation or the soft color method, represents an even more radical departure from the way color singlet states are conventionally
treated in perturbation theory. Color is, in fact, “ignored”. Rather than explicitly imposing that the system is in a color singlet state in the short-distance perturbative diagrams, the appearance of color singlet asymptotic states depends solely on the outcome of large-distance fluctuations of quarks and gluons. In other words, color is a nonperturbative phenomenon.

In Fig. 1 we show typical diagrams for the production of $\psi$-particles representing the competing treatments of the color quantum number. In the diagram of Fig. 1a, the color singlet approach, the $\psi$ is produced in gluon-gluon interactions in association with a final state gluon which is required by color conservation. This diagram is related by crossing to the hadronic decay $\psi \rightarrow 3$ gluons. In the color evaporation approach, the color singlet property of the $\psi$ is ignored at the perturbative stage of the calculation. The $\psi$ can, for instance, be produced to leading order by $q\bar{q}$-annihilation into $c\bar{c}$, which is the color-equivalent of the Drell-Yan process. This diagram is calculated perturbatively; its dynamics are dictated by short-distance interactions of range $\Delta x \simeq m_\psi^{-1}$. It does indeed not seem logical to enforce the color singlet property of the $\psi$ at short distances, given that there is an infinite time for soft gluons to readjust the color of the $c\bar{c}$ pair before it appears as an asymptotic $\psi$ or, alternatively, $D\bar{D}$ state. Alternatively, it is indeed hard to imagine that a color singlet state formed at a range $m_\psi^{-1}$, automatically survives to form a $\psi$. This formalism represents the original and, as we will show, correct method by which perturbative QCD calculations were performed.

![Typical diagrams for (a) color singlet $\psi$ production and (b) color evaporation $\psi$ production.](image)

Figure 1: Typical diagrams for (a) color singlet $\psi$ production and (b) color evaporation $\psi$ production.

We will first discuss the resolution of the charmonium problem, emphasizing the color evaporation approach. The solution suggests a radical departure from the way color is treated in perturbative QCD calculations. We will subsequently speculate on the implications for the dynamics underlying the production of rapidity gaps which refer to regions in phase space where no hadrons appears as a result of the production of a color neutral quark-antiquark pair. The connection to charmonium physics is obvious: the $\psi$ is a color-neutral $c\bar{c}$ pair!
2 Onium Calculations with Soft Color

The “color evaporation” or “soft color” treatment of the color quantum number lead to a similar description of bound and open charm production:

\[ \sigma_{\text{onium}} = \frac{1}{9} \int_{2m_c}^{2m_D} dm \frac{d\sigma_{c\bar{c}}}{dm}, \quad (1) \]

and

\[ \sigma_{\text{open}} = \frac{8}{9} \int_{2m_c}^{2m_D} dm \frac{d\sigma_{c\bar{c}}}{dm} + \int_{2m_c}^{m_D} dm \frac{d\sigma_{c\bar{c}}}{dm} \quad (2) \]

\[ \simeq \frac{8}{9} \int_{2m_c}^{2m_D} dm \frac{d\sigma_{c\bar{c}}}{dm} \quad (3) \]

where the cross section for producing heavy quarks, \( \sigma_{c\bar{c}} \), is computed perturbatively. Diagrams are included order-by-order, irrespective of the color of the \( c\bar{c} \) pair. The coefficients \( \frac{1}{9} \) and \( \frac{8}{9} \) represent the statistical probabilities that the \( 3 \times 3 \) charm pair is asymptotically in a singlet or octet state. In order to achieve the phenomenological success described here it is essential to systematically include next-to-leading order terms. Neglecting \( O(\alpha_s^3) \) terms is equivalent to describing photon interactions with matter neglecting the Bethe-Heitler process versus Compton scattering because it is a higher order process. The former actually dominates at high energy for reasons that are similar to those requiring the inclusion of higher order heavy quark processes.

In principal the calculation only predicts the sum of the cross sections of all onium states given by Eq. (1). This sum rule is, unfortunately, difficult to test experimentally, since it requires measuring cross sections for all of the bound states at a given energy. This does not mean that the calculation has no predictive power. The above equations make the bold prediction that all onium states \( \psi, \psi' \), \( \chi \) and \( \eta_c \) states share the same production dynamics which they also share with open charm in the limit \( m_c \simeq m_D \); see Eq. (3). The CDF collaboration has accumulated large samples of data on the production of prompt \( \psi, \chi_{cJ} \), and \( \psi' \). Since all charmonium states share the same production dynamics in the color evaporation scheme, their \( p_T \) distributions should be the same, up to a multiplicative constant. This prediction is borne out by the CDF data, as we can see in Fig. 2. We will return to a detailed calculation of the distribution shown further on. The formalism predicts furthermore that, up to color and normalization factors, the energy, \( x_F \) and \( p_T \) dependence of production cross sections for onium states and open charm pairs is the same. Support for the prediction of Eqs. (1) and (3) that the production of hidden and open charm have similar dynamics is shown in Fig. 3, which displays...
Figure 2: Data from the CDF Collaboration shown with arbitrary normalization. The curves are the predictions of the color evaporation model at tree level, also shown with arbitrary normalization. The normalization is correctly predicted within a K-factor of 2.2.

charm photoproduction data for both open charm and bound state production with common normalization in order to show their identical energy behavior. A similar figure for hadroproduction can be found in Ref. 11. By the same argument the formalism also predicts that the normalized $x_F$ distribution for $J/\psi$ and $D\bar{D}$ pairs should be the same at a given center-of-mass energy. This is indeed the case; see Fig.
Figure 3: Photoproduction data \[9,10\] as a function of the photon energy in the hadron rest frame, \(W_\gamma\). The normalization has been adjusted to show the similar shapes of the data.

One of the most striking features of color evaporation is that the production of charmonium is dominated by the conversion of a colored gluon into a \(\psi\), as in Fig. 1b. In the conventional treatment, where color singlet states are formed at the perturbative level, 3 gluons (or 2 gluons and a photon) are required to produce a \(\psi\). Contrary to the usual folklore, \(\psi\)'s are, except at the higher energies, not produced by gluons. As a consequence color evaporation predicts an enhanced \(\psi\) cross section for antiproton beams, while the color singlet model predicts roughly equal cross sections for proton and antiproton beams. The prediction of an enhanced \(\bar{p}\) yield is obviously correct: antiproton production of \(\psi\)'s exceeds that by protons by a factor 5 close to threshold; see Fig. 5. This fact has been known for some time.\[5,6,7\] We should note that for sufficiently high energies, gluon initial states will eventually dominate because they represent the bulk of soft partons.

3 Quantitative Tests of Soft Color

The color evaporation scheme assumes a factorization of the production of the \(c\bar{c}\) pair, which is perturbative and process dependent, and the materialization of this pair into a charmonium state by a mechanism that is nonperturbative and process independent. This assumption is reasonable since the characteristic time scales of the two processes are very different: the time scale for the production of the pair is the inverse of the heavy quark mass, while the formation of the bound state is longer than the time scale \(1/\Lambda_{QCD}\). Therefore, explicit comparison with the \(\psi\) data requires knowledge of the fraction \(\rho_\psi\) of
produced onium states that materialize as $\psi$'s, i.e.,

$$\sigma_\psi = \rho_\psi \sigma_\text{onium} ,$$

where $\rho_\psi$ is assumed to be a constant, independent of the process. This assumption is in agreement with the low energy data. The constant not only accounts for the direct production of $\psi$ to the onium cross section, but also includes its production via $\psi'$ and $\chi$ production and decay.

Quantitative tests of color evaporation are made possible by the fact that all $\psi$-production data, i.e. photo-, hadroproduction, $Z$-decay, etc., are described in terms of a single parameter. Once $\rho_\psi$ has been empirically determined for one initial state, the cross section is predicted without free parameters for the other. We will illustrate the power of the color evaporation scheme by showing how it quantitatively accommodates all measurements, including the high energy Tevatron and HERA data, which have represented a considerable challenge for the color singlet model. Its parameter-free prediction for the $Z$-boson decay rate into $\psi$'s is an order of magnitude larger than the color singlet model and consistent with data.

In Fig. 6 we compare the photoproduction data with theory, using the NLO perturbative QCD calculation of charm pair production from Ref. 17. From the relative magnitude of the $\psi$ and open charm cross sections we determine the fragmentation factor $\rho_\psi$ to be 0.50 using GRV HO, or 0.43 using MRS A structure functions. Note that the factor $\rho_\psi$ possesses a theoretical uncertainty due to the choice of scales and parton distribution functions. We conclude the
photoproduction of $J/\psi$ and $D\bar{D}$ is well described by the color evaporation model. This reaction has now been used to fix the only free parameter, $\rho_\psi \approx 0.5$.

At this point the predictions of the color evaporation model for hadroproduction of $\psi$ are completely determined, up to $O(\alpha_s^4)$ QCD corrections. In Fig. 7 we compare the color evaporation model predictions with the data and conclude that the this color scheme describes the hadroproduction very accurately. In order to also obtain a theoretical prediction for the $p_T$-distribution already shown in Fig. 2, we have computed the processes $g + g \rightarrow [c\bar{c}] + g$, $q + q \rightarrow [c\bar{c}] + g$, and $g + g \rightarrow [c\bar{c}] + q$ at tree level using MADGRAPH. We imposed that the $c\bar{c}$ pair satisfy the invariant mass constraint of Eq. (1).

Our results are shown in Fig. 8. Higher order corrections such as soft-gluon resummation are expected to tilt our lowest order prediction, bringing it to a closer agreement with the data.

In the color-evaporation scheme the width for inclusive $Z$ decay into prompt
Figure 6: Photoproduction data\cite{9,10} and the predictions of the color evaporation model at next-to-leading order as a function of the photon energy in the hadron rest frame, $W_\gamma$. The normalizations in this figure are absolute.

The normalizations in this figure are absolute.

charmonium is:

$$\Gamma(Z \rightarrow \text{prompt charmonium}) = \frac{1}{9} \int_{2m_c}^{2m_D} dm \frac{d\hat{\Gamma}_{c\bar{c}}}{dm},$$

where $\hat{\Gamma}$ is the partonic width for producing a $c\bar{c}$ pair. The procedure should be familiar: in order to obtain the partial width into a specific charmonium state we multiply the above expression by the appropriate fragmentation fraction $\rho$ into $\psi$, which was determined from charmonium photoproduction data. Notice that the predictions for the $Z$ decay into charmonium are parameter-free.

We have again evaluated all the tree-level partonic amplitudes using the package MADGRAPH\cite{18}. Although formally of higher order in $\alpha_s$, the dominant process for the inclusive decay of the $Z$ into charmonium is $Z \rightarrow c\bar{c}q\bar{q}$, where $q = u, d, s, c, b$. (The leading-order process in $\alpha_s$ is $Z \rightarrow c\bar{c}g$, which leads to the production of a charmonium state and a hard jet is suppressed by a virtual quark propagator of order $m_c/m_Z$). The branching fraction of $Z$ into prompt $\psi$ is $(1.7–1.8) \times 10^{-4}$. This is to be contrasted with the color-singlet model which predicts a branching fraction for direct $\psi$ in $Z$ decay of the order
Figure 7: Hadroproduction data\textsuperscript{20,21} and the predictions of the color evaporation model at next-to-leading order as a function of the center-of-mass energy, $E_{\text{cm}}$. The curve for bound state production is an absolutely normalized, parameter-free prediction of the color evaporation model.
The color-evaporation model leads to a branching fraction larger by almost an order of magnitude consistent with the result reported by the OPAL collaboration of

\[ B(Z \to \text{prompt } \psi + X) = (1.9 \pm 0.7 \pm 0.5 \pm 0.5) \times 10^{-4} . \]

We hope that we have illustrated by now that the soft color approach gives a complete picture of charmonium production in hadron-hadron, \( \gamma \)-hadron, and \( Z \) decays. The phenomenological success of the soft color scheme is impressive and extends to applications to other charmonium and upsilon states.

### 4 Intermezzo: Soft Color and BBL

Other approaches, very similar in spirit, can be found in Refs. 3, 15 and 23. The color evaporation approach differs from Ref. 3, the formalism of Bodwin, Braaten and Lepage, in the way that the \( c\bar{c} \) pair exchanges color with the underlying event. In the BBL formalism, multiple gluon interactions with the \( c\bar{c} \) pair are suppressed by powers of \( v \), the relative velocity of the heavy quarks within the \( \psi \). The color evaporation model assumes that these low-energy interactions can take place through multiple, soft-gluon interactions. While the formalism allows straightforward application to heavy quark decays, it is not always clear how to compute production cross sections in the BBL formalism (e.g. photoproduction of \( \psi \) near \( z = 1 \)). The color evaporation scheme, though partly nonperturbative, is phenomenologically well-defined, has less parameters (1 versus 3 for describing \( \psi \)-production). Also, next-to-leading order corrections are included in a straightforward way, a necessary condition for obtaining quantitative predictions.

The \( \psi \)'s produced through the color-evaporation mechanism are expected to be unpolarized since the polarization information is lost because of the multiple soft gluon exchanges. On the other hand, the (non)polarization of \( \psi \) is hard to explain in the framework of the color-octet model. Therefore, the measurement of the polarization of the produced charmonium may very well be a tool to discriminate between these competing descriptions.

### 5 Implications for the Physics of Rapidity Gaps

The important lesson about color resides however in the similarity, not the differences of these approaches: perturbative color octet states fully contribute to the asymptotic production of color singlet states such as \( \psi \)'s. We suspect that this is also true for the production of a rapidity gap which is, e.g. when produced in electroproduction, nothing but the creation of a color singlet.
quark-antiquark pair; see Fig. 8. The diagram shown represents the production of final state hadrons which are ordered in rapidity. From top to bottom we find the fragments of the intermediate partonic quark-antiquark state and those of the target. Buchmüller and Hebecker proposed that the origin of a rapidity gap corresponds to the absence of color between photon and proton, i.e. the $3 \times \bar{3} (= 1 + 8)$ intermediate quark-antiquark state is in a color singlet state. Because color is the source of hadrons, only the color octet states yield hadronic asymptotic states. This leads to the approximate expectation that

$$F_2^{(gap)} = \frac{1}{1 + 8} F_2$$

Although this result is subject to corrections, it embodies the essential physics: events with and without gaps are described by the same short-distance dynamics. Essentially non-perturbative final-state interactions dictate the appearance of gaps whose frequency is determined by simple counting. The treatment of color is the same as in the case of heavy quark production: the same perturbative mechanisms, i.e. gluon exchange, dictates the dynamics of color-singlet gap ($\psi$) and regular deep inelastic (open charm) events.

Our understanding of the (soft) nature of color challenges the orthodox description of rapidity gaps in terms of the so-called hard Pomeron description sketched in Fig. 9. The $t$-channel exchange of a pair of gluons in a color singlet state is the origin of the gap. The color string which connects photon and proton in diagrams such as the one in Fig. 8 is absent and no hadrons are produced in the rapidity region separating them. The same mechanism predicts rapidity gaps between a pair of jets produced in hadronic collisions; see Fig. 10. These have been observed and occur with a frequency of order of one percent. The arguments developed in this work invalidate this approach: it is as meaningless to enforce the color singlet nature of the gluon pair as it is to require that the $c\bar{c}$ pair producing $\psi$ is colorless at the perturbative...
Following our color scheme the gaps are accommodated as a mere final state color bleaching phenomenon \textit{à la} Buchmüller and Hebecker. This can be visualized using the diagram shown in Fig. 11. At short distances it represents a conventional perturbative diagram for the production of a pair of jets. Also shown is the string picture for the formation of the final state hadrons. Color in the final state is bleached by strings connecting the $3$ jet at the top with the $3$ spectator di-quark at the bottom and vice-versa. The probability to form a gap can be counted \textit{à la} Buchmüller and Hebecker to be $1/(1+8)^2$ because it requires the formation of singlets in 2 strings. This is consistent with observation and predicts that, as was the case for electroproduction, the same short distance dynamics governs events with and without rapidity gaps. The data is consistent with the prediction of this simple picture which basically predicts that the gap fraction between $pp$ jets is the square of that between virtual photon and proton in deep inelastic scattering.

One should realize that this string picture is not necessarily the correct one. It is more likely that the color is bleached between the top and bottom $3$
Figure 11: Color bleaching picture for the formation of rapidity gaps in hadron collisions.

and $\bar{3}$ which are widely separated in rapidity space.

This discussion ignores that gluon-gluon as well as quark-quark subprocesses contribute to jet production in hadron collisions. In the color flow diagram corresponding to Fig. 11 top and bottom protons each split into a color octet gluon and color octet 3-quark remnant. There are now $(8 \times 8)^2$ color final states. We anticipate a reduced probability to form a color singlet. The reduction may not be very significant because the $10 + \bar{10}$ and 27 color final states may be suppressed. One argument for this is that these representations consist of exotic multi-quark states which do not materialize into final state mesons. High color charges may also be suppressed for dynamical reasons. Despite the fact that we can at best guess the non-perturbative dynamics, it is clear that the soft color formalism predicts a gap rate which is similar of smaller in gluon-gluon interactions. This is in contrast with the diagram of Fig. 9 which predicts a gap-rate enhanced by a factor $(\frac{4}{3})^2$ in gluon-gluon subprocesses. The contrasting predictions can be easily tested by enhancing the relative importance of quark-quark subprocesses, i.e. by increasing the $p_T$ of the jets at fixed energy of by decreasing the collision energy of the hadrons at fixed $p_T$. In either case we anticipate in the soft color scheme an increased rate for the production of gaps, a prediction opposite from that expected in the 2-gluon exchange model.

Do Figs. 8-11 suggest that we have formulated alternative $s$- and $t$-channel pictures to view the same physics? Although they seem at first radically different, this may not be the case. Computation of the exchange of a pair of colorless gluons in the $t$-channel is not straightforward and embodies all the unsolved mysteries of constructing the “Pomeron” in QCD. In a class of models where the Pomeron is constructed out of gluons with a dynamically generated mass, the diagram of Fig. 10 is, not surprisingly, dominated by the configuration where one gluon is hard and the other soft. The diagram
is identical to the standard perturbative diagram except for the presence of a soft, long-wavelength gluon whose only role is to bleach color. Its dynamical role is minimal, events with gaps are not really different from events without them. Soft gluons readjust the color at large distances and long times. Their description is outside the realm of perturbative QCD. In this class of models the hard Pomeron is expected to be no more than an order $\alpha_s^2$ correction, a view which can be defended on more solid theoretical ground. Some have challenged the theoretical soundness of this line of thinking. Also note that our discussion is at best indirectly relevant to completely non-perturbative phenomena like elastic scattering. There is no short distance limit defined by a large scale. The Pomeron exists.

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1. See, e.g., E. Braaten, S. Fleming, and T.C. Yuan, preprint OHSTPY-HEP-T-96-001 [hep-ph/9602374], to appear in Ann. Rev. Nucl. Part. Sci., and references therein.
2. W. Buchmüller, Phys. Lett. B353, 335 (1995); W. Buchmüller and A. Hebecker Phys. Lett. B355, 573 (1995).
3. G.T. Bodwin, E. Braaten, and G. Lepage, Phys. Rev. D51, 1125 (1995).
4. H. Fritzsch, Phys. Lett. B67, 217 (1977).
5. F. Halzen, Phys. Lett. B69, 105 (1977).
6. F. Halzen and S. Matsuda, Phys. Rev. D17, 1344 (1978).
7. M. Gluck, J. Owens, and E. Reya, Phys. Rev. D17, 2324 (1978).
8. CDF Collaboration, F. Abe, et al., preprint FERMILAB-CONF-95/128-E; preprint FERMILAB-CONF-95/136-E; preprint FERMILAB-CONF-95/263-E.
9. P.L. Frabetti, et al., Phys. Lett. B316, 197 (1993); R. Barate, et al., Z. Phys. C33, 505 (1987); M. Binkley, et al., Phys. Rev. Lett. 48, 73 (1982); M. Arneodo, et al., Phys. Lett. B332, 195 (1994); B. H. Denby, et al., Phys. Rev. Lett. 52, 795 (1984); M. Derrick, et al., Phys. Lett. B350, 120 (1995).
10. M.I. Adamovich et al., Phys. Lett. B187, 437 (1987); M.P. Alvarez et al., Z. Phys. C60, 53 (1993); K. Abe et al., Phys. Rev. D33, 1 (1986); J.J. Aubert et al., Nucl. Phys. B213, 31 (1983); J.C. Anjos et al., Phys. Rev. Lett. 62, 513 (1989); M. Derrick, et al., Phys. Lett. B349, 225 (1995).

11. J. Amundson, O. Éboli, E. Gregores, and F. Halzen, Phys. Lett. B372, 127 (1996).

12. LEBC-EHS Collaboration, M. Aguilar-Benitez et al., Z. Phys. C40, 321 (1988).

13. E705 Collaboration, L. Antoniazzi, et al., Phys. Rev. D46, 4828 (1992).

14. WA39 Collaboration, M.J. Corden, et al., Phys. Lett. B98, 220 (1981); NA3 Collaboration, J. Badier, et al., Z. Phys. C20, 101 (1983); K.J. Anderson, et al., Phys. Rev. Lett. 42, 944 (1979); UA6 Collaboration, C. Morel, et al., Phys. Lett. B252, 505 (1990); M.J. Corden, et al., Phys. Lett. B68, 96 (1977).

15. R. Gavai, et al., Int. J. Mod. Phys. A10, 3043 (1995).

16. G.A. Schuler, preprint CERN-TH.7170/94 [hep-ph/9403387].

17. R.K. Ellis and P. Nason, Nucl. Phys. B312, 551 (1989); P. Nason, S. Dawson, and R.K. Ellis, ibid. B327, 49 (1989).

18. W. Long and T. Steltzer, Comput. Phys. Commun. 81, 357 (1994).

19. S. Fleming, private communication and J. Amundson and S. Fleming, preprint MADPH-95-914 [hep-ph/9601298].

20. A. Bamberger et al., Nucl. Phys. B134, 1 (1978); J.J. Aubert et al., Phys. Rev. Lett. 33, 1404 (1974); WA39 Collaboration, M.J. Corden et al., Phys. Lett. B98, 220 (1981); Yu.M. Antipov et al., ibid. B60, 309 (1976); K.J. Anderson et al., Phys. Rev. Lett. 37, 799 (1976); NA3 Collaboration, J. Badier et al., Z. Phys. C20, 101 (1983); K.J. Anderson et al., Phys. Rev. Lett. 42, 944 (1979); E705 Collaboration, L. Antoniazzi et al., ibid. 70, 383 (1993); UA6 Collaboration, C. Morel et al., Phys. Lett. B252, 505 (1990); H.D. Snyder et al., Phys. Rev. Lett. 36, 1415 (1976); E672/706 Collaboration, V. Abramov et al., FERMILAB-PUB-91-62-E, March 1991; E789 Collaboration, M.H. Schub et al., Phys. Rev. D52, 1307 (1995); E. Nagy et al., Phys. Lett. B60, 96 (1975).

21. H. Cobbaert et al., Z. Phys. C36, 577 (1987); ACCMOR Collaboration, S. Barlag et al., ibid. C39, 451 (1988); J.L. Ritchie et al., Phys. Lett. B126, 499 (1983); LEBC-EHS Collaboration, M. Aguilar-Benitez et al., ibid. B122, 312 (1983); B123, 103 (1983); LEBC-MPS Collaboration,
R. Ammar et al., Phys. Lett. B183, 110 (1987), erratum B192, 478 (1987); R. Ammar et al., Phys. Rev. Lett. 61, 2185 (1988); A.G. Clark et al., Phys. Lett. B77, 339 (1978).

22. V. Barger, K. Cheung, and W.-Y. Keung, Phys. Rev. D41, 1541 (1990).

23. P. Hoyer and C. Lam, preprint NORDITA-95/53 P [hep-ph/9507367].

24. A. Brandenburg, O. Nachtmann, and E. Mirkes, Z. Phys. C60, 697 (1993).

25. M. Beneke and I.Z. Rothstein, SLAC-PUB-7129 [hep-ph/9603400]; W.K. Tang and M. Vänttinen, Phys. Rev. D53, 4851 (1996); NORDITA-96-18-P [hep-ph/9603266].

26. E. Braaten and Yu-Qi Chen, OHSTPY-HEP-T-96-010 [hep-ph/9604237].

27. D0 Collaboration, S. Abachi et al., preprint FERMILAB-PUB-95/302-E.

28. D. Zeppenfeld, preprint MADPH-95-933 [hep-ph/9603315].

29. F. Halzen, G. Ingelman and D. Zeppenfeld, private communication.

30. A. A. Natale, et al. Phys. Rev. D47 (1993) 295; D48 2324 (1993).

31. H. Chehime, et al., Phys. Lett. B286 397 (1992).

32. J. R. Cudell, A. Donnaichie, and P. V. Landshoff, to appear.

33. J.D. Bjorken, Int. J. Mod. Phys. A7 4189 (1992); Phys. Rev. D47 101 (1993); preprint SLAC-PUB-5823 (1992).

34. A. White, talk given at the 3rd Workshop on Small x and Diffractive Physics, Argonne, Sept. 26–29, 1996.