A THEORETICAL REVIEW OF THE PHOTON STRUCTURE FUNCTION

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I. Introduction

We are seeing at this workshop the results of almost incredible advances on the experimental side of two photon physics. We are also lucky, perhaps, that there have been exciting and important new advances on the theoretical side as well. The main improvements are that the reasons for the somewhat bothersome higher order predictions of negative values of the structure function $F_2^{\gamma}(x,Q^2)$ are now more thoroughly understood, and that the range of applicability of the QCD predictions for F_2^{γ} is more carefully delineated and appreciated. The implications of this new understanding for the old experimental goal of using F_2^{γ} to measure Λ are not, as we shall discuss, altogether positive.

II. A Brief History

The prediction and experimental discovery of approximate Bjorken scaling in deep inelastic electron nucleon scattering has taught us to think in terms of pointlike constituents of hadrons dominating the physics, at least in certain kinematic regions. In 1971 the extension of these ideas to two photon physics was suggested by Walsh[1)] and by Brodsky, Kinoshita and Terazawa[2)]. These authors suggested that the reaction

$$e^+ e^- \to e^+ e^- + \text{hadrons}$$

would offer an opportunity to measure the structure functions of (nearly) real target photons when probed by another highly virtual photon by tagging one lepton at a large angle and the other lepton at a small angle. When the hadronic properties of the target photon are interpreted in terms of the vector meson dominance (VMD) model, then one expects to observe approximate Bjorken scaling for the photon structure function, exactly in analogy with the nucleon target case.

In 1973 Walsh and Zerwas[3)] pointed out that the target photon could actually couple directly to its pointlike quark-parton constituents and provide anomalous contributions, e.g. $F_2^\gamma(x,Q^2)$ has a nonscaling contribution proportional to ln Q^2/m^2, and $F_L^\gamma(x,Q^2)$ is nonzero in distinction to the usual expectations for spin 1/2 partons. This subject was later studied by Kingsley[4)] and by Worden.[5)]

In 1974 it was discovered by Politzer[6)] and by Gross and Wilczek[7)] that nonabelian gauge theories (e.g. QCD) are asymptotically free and offer an explanation of the approximate Bjorken scaling observed in deep inelastic electron nucleon scattering, as well as quantitative predictions for the pattern of scaling violations to be expected. The predictions of QCD for the hadronic or VMD part of the photon structure functions were then studied by Ahmed and Ross[8)] in 1975.

The treatment of the photon structure functions in QCD was the subject of a remarkable paper by Witten[9)] in 1977. Witten showed, within the operator product expansion (OPE), how the hadronic and pointlike parts of the photon structure functions are unified by operator mixing between the usual quark and gluon operators and the photon operator. The main new result was that the photon structure function is absolutely calculable in QCD for asymptotically large values of Q^2.

In 1979 Witten's OPE analysis was translated into diagrammatic and/or Altarelli-Parisi-Lipatov[10)] language by Llewellyn Smith[11)], Frazer and Gunion[12)], and DeWitt, Jones, Sullivan, Willen and Wyld.[13)]

The completion of the calculations of the higher order QCD corrections to deep inelastic scattering in 1978 led Bardeen and Buras[14)] in 1979 to extend Witten's OPE analysis beyond the leading order. In 1980 Duke and Owens[15)] extended slightly the Bardeen-Buras calculation and pointed out that the higher order corrections actually drove the perturbatively calculable part of F_2^γ to negative values for $x<0.1$ and $x\sim 1$.

For experimentally accessible values of Q^2, these unphysical predictions of negative cross sections are not alleviated by simply adding the hadronic VMD pieces in the usual naive way.

In 1981 Uematsu and Walsh[16] showed that when the target photon mass $-p^2$ is not zero and bounded by $\Lambda^2 << p^2 << Q^2$, then the hadronic part of F_2^γ is perturbatively calculable and the predictions, including higher order corrections, are no longer negative for small x (although there is still a problem for large $x\sim 1$).

Based on the Uematsu-Walsh result, Bardeen[17] explained the origin and resolution of the small x problem in his 1981 Bonn Conference talk. In addition, Frazer[18] has discussed the behavior of $F_2^\gamma(x,Q^2)$ for $x\sim 1$ based on Frazer and Rossi[19], and Chase.[20]

III. Witten's Breakthrough

Prior to the work of Witten, the photon structure function was generally understood in terms of two separate pieces - a hadronic piece presumably described by VMD, and a pointlike piece described by the box diagram. Neither of these pieces is perturbatively calculable in QCD, since we cannot calculate either the hadronic matrix elements of the photon for VMD or the dynamical mass scale(s) necessary to regulate the infrared divergences of the box diagram. In terms of moments of $F_2^\gamma(x,Q^2)$, for example, we have then

$$M_n^\gamma(Q^2) = \int_0^1 dx\, x^{n-2}\, F_2^\gamma(x,Q^2)$$
$$= \sum_{i=+,-,NS} C_n^i(Q^2/\ ^2, g^2,\ )\ \langle\gamma|O_n^i|\gamma\rangle$$
$$+ M_n^{box}(Q^2). \qquad (1)$$

In Eq.(1) the C_n^i are each of $O(\alpha^0)$ while the $\langle\gamma|O_n^i|\gamma\rangle$ are each of $O(\alpha^1)$. Also, to $O(\alpha)$ we have

$$M_n^{box}(Q^2) = \frac{3f\langle e^4\rangle\alpha}{\pi} \times \left[(1-2x+2x^2)\ \ln\frac{Q^2(1-x)}{m^2x} -1+8x+8x^2\right].$$

Witten realized that in the OPE framework, we should have instead

$$M_n^{box}(Q^2) = C_n^{\gamma}(Q^2/\mu^2, g^2, \alpha)\langle\gamma|O_n^{\gamma}|\gamma\rangle, \tag{2}$$

where now $C_n^{\gamma}\sim O(\alpha)$ and $\langle\gamma|O_n^{\gamma}|\gamma\rangle\sim O(1)$. In Eq.(2) we note the appearance of a twist-two operator O_n^{γ} constructed from the photon tensor $F_{\mu\nu}$ in the same way that O_n^G is consructed from the gluon tensor $G_{\mu\nu}$.

The essential new idea - mixing between the operators $O_n^{+,-,NS}$ and O_n^{γ} - is expressed by the anomalous dimension matrix[9]

$$\begin{bmatrix} \gamma_n^{++} & \gamma_n^{--} & 0 & K_n^{+} \\ \gamma_n^{-+} & \gamma_n^{--} & 0 & K_n^{-} \\ 0 & 0 & \gamma_n^{NS} & K_n^{NS} \\ 0 & 0 & 0 & 0 \end{bmatrix}.$$

Applying the renormalization group treatment to $M_n^{\gamma}(Q^2)$ in the usual way, we get[9,14]

$$\begin{aligned} M_n^{\gamma}(Q^2) = &\sum_{i,j} C_n^i(1,\bar{g}^2,\alpha)\ M_n^{i,j}\ \langle\gamma|O_n^j|\gamma\rangle \\ &+ \sum_i C_n^i(1,\bar{g}^2,\alpha)\ X_n^i\ \langle\gamma|O_n^{\gamma}|\gamma\rangle \\ &+ C_n^{\gamma}(1,\bar{g}^2,\alpha)\ \langle\gamma|O_n^{\gamma}|\gamma\rangle, \quad i,j=+,-,NS \end{aligned} \tag{3}$$

Now using $\langle\gamma|O_n^{\gamma}|\gamma\rangle=1$ to $O(\alpha^0)$, we finally get

$$\begin{aligned} M_n^{\gamma}(Q^2) = &\sum_i A_n^i(\mu^2)\left(\frac{\alpha_s(Q^2)}{\alpha_s(\mu^2)}\right)^{d_n^i} \\ &+\sum_i \frac{a_n^i}{\alpha_s(Q^2)}\ \frac{1}{d_n^i+1}\left[1-\left(\frac{\alpha_s(Q^2)}{\alpha_s(\mu^2)}\right)^{d_n^i+1}\right] \\ &+\sum_i \frac{b_n^i}{d_n^i}\left[1-\left(\frac{\alpha_s(Q^2)}{\alpha_s(\mu^2)}\right)^{d_n^i}\right] + c_n^{\gamma} \end{aligned} \tag{4}$$

In Eq.(4) μ^2 is an arbitrary renormalization point, the A_n^i are unknown hadronic matrix elements of the photon, the d_n^i are proportional to the one-loop anomalous dimensions (e.g. $d_n^{NS} = \gamma_n^{NS}/2\beta_0$) , and the a_n^i, b_n^i are exactly calculable (indeed, calculated) numbers.

Now for Q^2 asymptotically large, we have $\alpha_s(Q^2) \sim 1/\ln Q^2$ so for fixed n we see that $1/\alpha_s$ dominates $(\alpha_s(Q^2)/\alpha_s(\mu^2))$ if $Q^2 >> \Lambda^2$ and we have

$$\lim_{\substack{Q^2 \to \infty \\ n \text{ fixed}}} M_n^\gamma(Q^2) = \frac{1}{\alpha_s(Q^2)} \sum_i \frac{a_n^i}{d_n^i+1} + \sum_i \frac{b_n^i}{d_n^i} + c_n^\gamma \tag{5}$$

Eq.(5) is the source of the often heard statement: " QCD makes an absolute, one parameter(Λ) prediction for $F_2^\gamma(x,Q^2)$". The shape of $F_2^\gamma(x,Q^2)$ implied by Eq.(5) is noticeably different from that due to the box graph alone (see fig. 1). This change is due to the renormalization or mixing effects of the hadronic operators. Note that the leading logarithm (LL) result of Witten is much smaller than the LL part of the box at large x, and that the LL result of Witten has developed a small peak at $x \sim 0$. This overall behavior is easily understood by separating in the calculation the parts proportional to $\langle e^4 \rangle$ and $\langle e^2 \rangle^2$, respectively.[15] The former are usually termed "valence" contributions because the struck quark originates directly from the parent target photon (see figs. 2,3). The latter are called "sea" contributions because the struck quark was generated by evolution, or radiation from the valence quarks. The parallel with the nucleon target case is then rather obvious.

IV. Higher Order Breakdown and Its Resolution

As we have just discussed, the asymptotic predictions for $M_n^\gamma(Q^2)$ are calculable with only the QCD scale Λ needed in order to make an absolute prediction. It came as somewhat of a discouraging surprise, however, when it was found[15] that, taken as it stands, the result

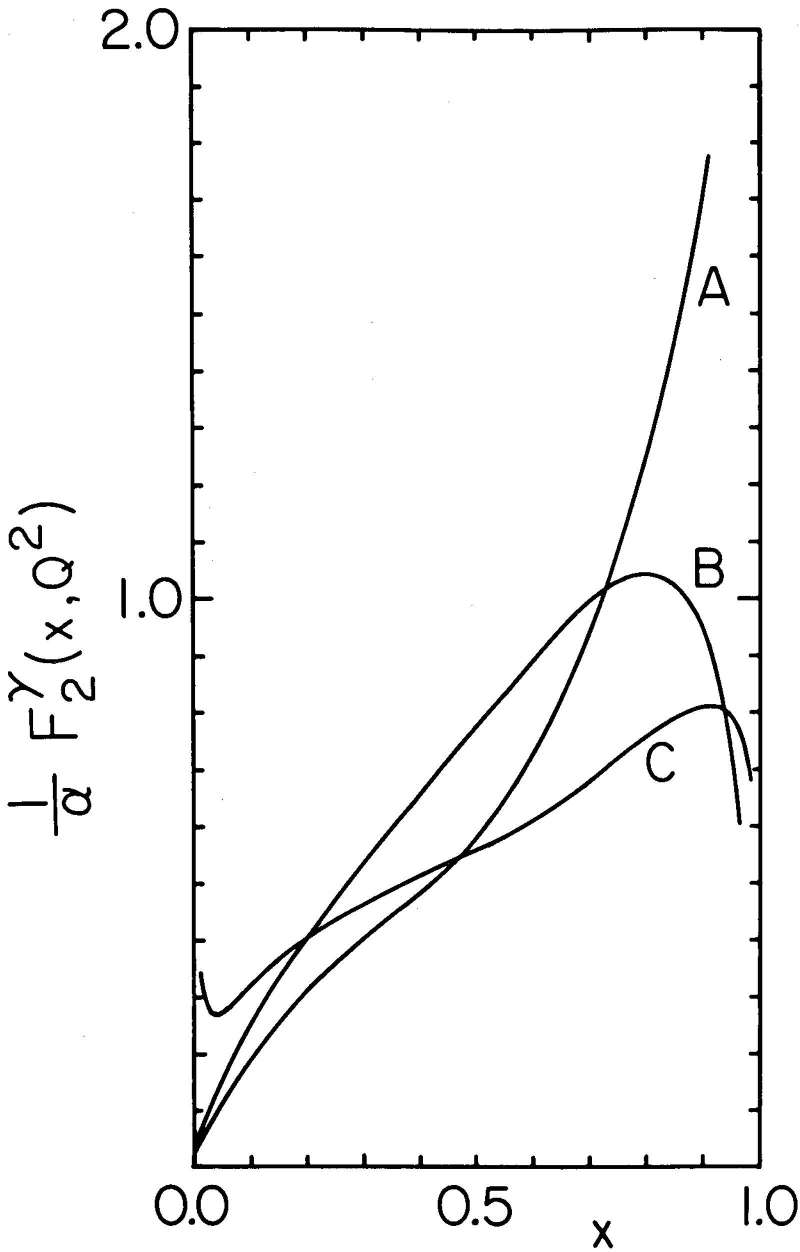

Fig. 1. A) The leading log part of the box graph.
B) The full box graph contribution.
C) Witten's leading log result.

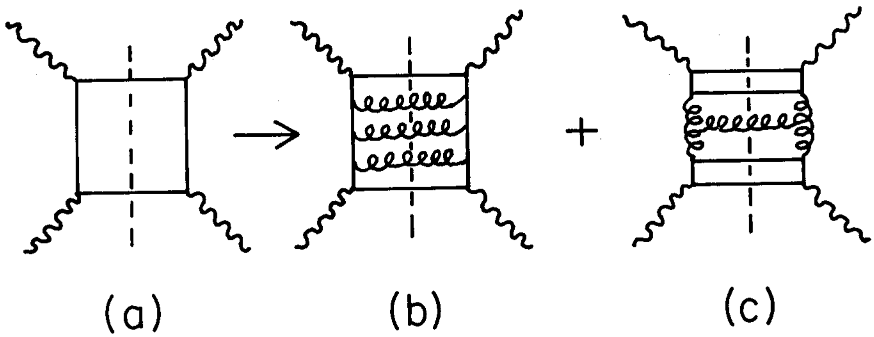

Fig. 2. A schematic demonstration of how QCD evolution turns the box graph(a) into a valence part(b) and a sea part(c).

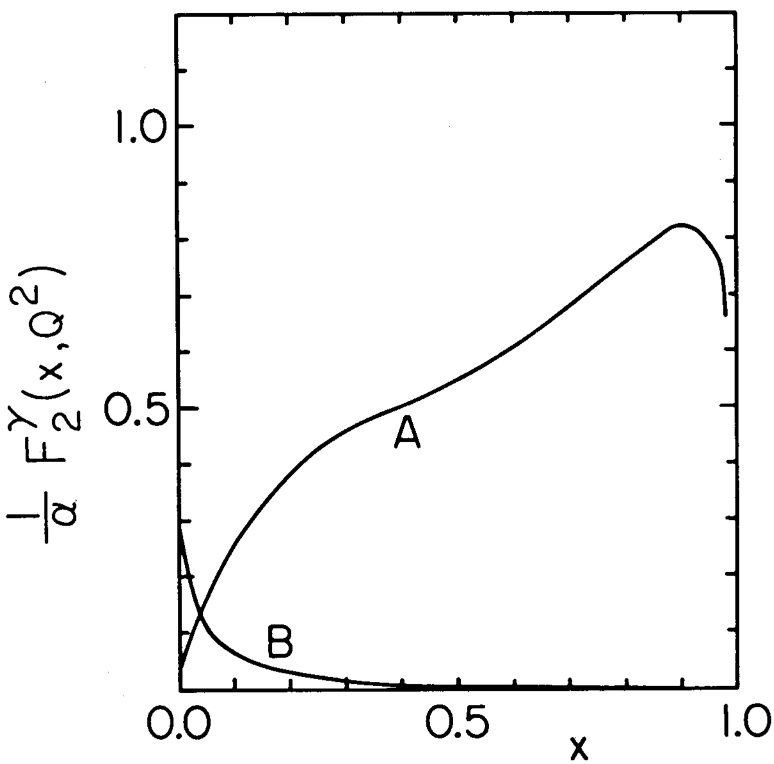

Fig. 3. A) The valence part of Witten's leading log result.
B) The corresponding sea part.

Eq.(5) including the higher order corrections is really quite sick for $x\sim0$, $x\sim1$ where it predicts negative values of $F_2^{\gamma}(x,Q^2)$ (see fig. 4). For $x\sim0$ the damage is caused by the $\langle e^2\rangle^2$ "sea" term, which has a large negative spike. For $x\sim1$ the problem is in the $\langle e^4\rangle$ term whose moments are actually negative for large enough n.

In order to understand the origin of these problems and hence their resolution, we need only look a bit more carefully at the steps leading from Eq.(4) to Eq.(5) above.

A. $x\sim0$

In order to reconstruct $F_2^{\gamma}(x,Q^2)$ we must perform an inverse Mellin transform on $M^{\gamma}(n,q^2)$:

$$F_2^{\gamma}(x,Q^2) = \frac{1}{2\pi i}\int_{c-i\infty}^{c+i\infty} dn\ x^{-n+1}\ M^{\gamma}(n,Q^2) \qquad (6)$$

Now the rightmost singularity in n of $M^{\gamma}(n,Q^2)$ controls the leading $x\sim0$ behavior of $F_2^{\gamma}(x,Q^2)$, i.e. if

$$\lim_{n\to n_0} M^{\gamma}(n,Q^2) \sim \frac{c}{n-n_0} + \text{regular},$$

then Eq.(6) implies

$$\lim_{x\to 0} F_2^{\gamma}(x,Q^2) \sim cx^{1-n_0}.$$

Suppose we apply Eq.(6) to Eq.(5) including the next to leading terms b_n. Then the dominant pole in n occurs when $d_n^- = 0$ or $n_0 = 2$, leading to

$$F_2^{\gamma}(x,Q^2) \sim -c/x.$$

This is precisely the behavior observed for the sea term in fig. 4. In addition, the LL result of Witten has a pole when $d_n^- + 1 = 0$ or $n_0 = 1.5962$ leading to the positive spike $\sim x^{-0.5962}$ observed in the LL sea (see fig. 3).

It is clear, however, that these predictions of spikes are entirely spurious. To see this, look at Eq.(4) and note that the poles

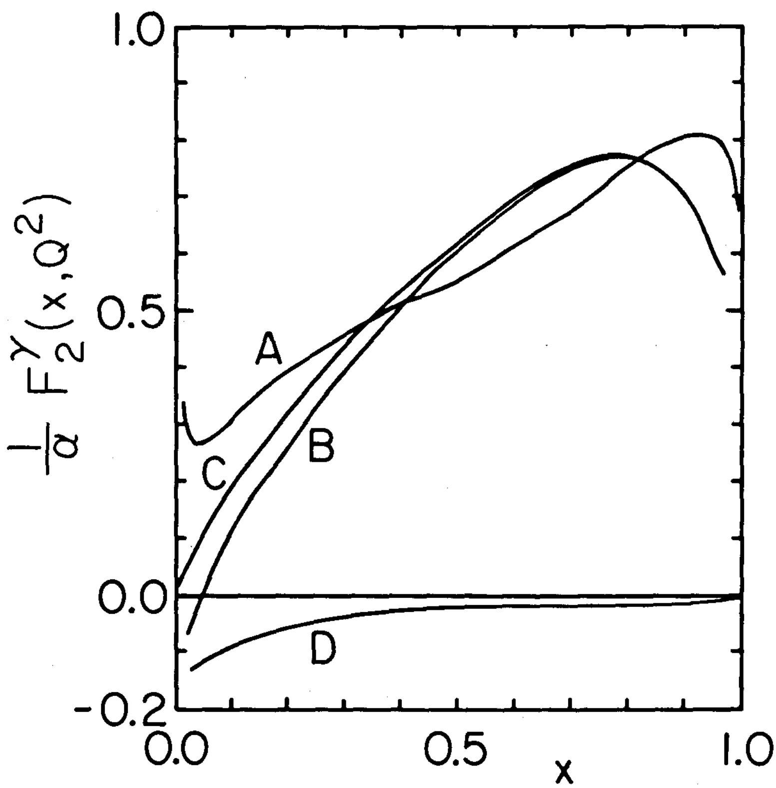

Fig. 4. The LL result of Witten A) compared to the higher order calculation of Bardeen and Buras B). Curve C) is the valence part and curve D) is the sea part of the higher order calculation.

$1/d$ are always accompanied by factors $(1 - x^d)$, where $x = \alpha_s(Q^2)/\alpha_s(\mu^2)$ and $d = d_n^i+1$, d_n^i,...etc. Instead of a pole we actually get for $d \to 0$,

$$\lim_{d\to 0} \frac{1}{d}(1 - x^d) = -\ln x.$$

This simply means that we must pay strict attention to the order of the limits $Q^2 \to \infty$, $n \to n_0$ - they are not interchangeable. Of course this kind of nonuniform convergence of the moments is well-known in the nucleon or hadronic sector[21); it is just that the penalty for ignoring it seems much more severe in the photon target case.

The behavior noted above was first observed[16) in the case of deep inelastic scattering on an off-shell photon of virtual mass $-p^2$. In that case even the hadronic matrix elements A_n^i are perturbatively calculable for $p^2 \gg \Lambda^2$ and it is natural then to choose $\mu^2 = p^2$ in Eq.(4). This leads to positive predictions for $F_2^\gamma(x,Q^2)$ for $x\sim 0$ but unfortunately we lose the ability to use F_2^γ to measure Λ! In effect the dominant logarithm becomes $\ln Q^2/p^2$ instead of $\ln Q^2/\Lambda^2$. This is illustrated in fig. 5.

Now for real photons[17) ($p^2 < \Lambda^2$) there is still some freedom in the method in which the $1/d_n$ poles are regulated. To see this note that Eq.(4) can be rewritten trivially as

$$M_n^\gamma(Q^2) = \sum_i \tilde{A}_n^i(\mu^2)\left(\frac{\alpha_s(Q^2)}{\alpha_s(\mu^2)}\right)^{d_n^i} + \frac{1}{\alpha_s(Q^2)}\sum_i \frac{a_n^i}{d_n^i+1} + \sum_i \frac{b_n^i}{d_n^i} + c_n^\gamma,$$

where the $\tilde{A}^i{}_n$ differ from the already unknown A_n^i only by a finite, and hence harmless, renormalization. Thus there are poles in the $\tilde{A}_n^i$ but for $p^2 < \Lambda^2$ we have no particular reason to assume the exact form of Eq.(7). For example, a set of simple poles $f(n)/n-n_0$ with an arbitrary $f(n)$ that gives the correct residue would be a possibility. A simple scheme along these lines has been worked out by Antoniadis and Grunberg[22). In their scheme there is one free parameter t which

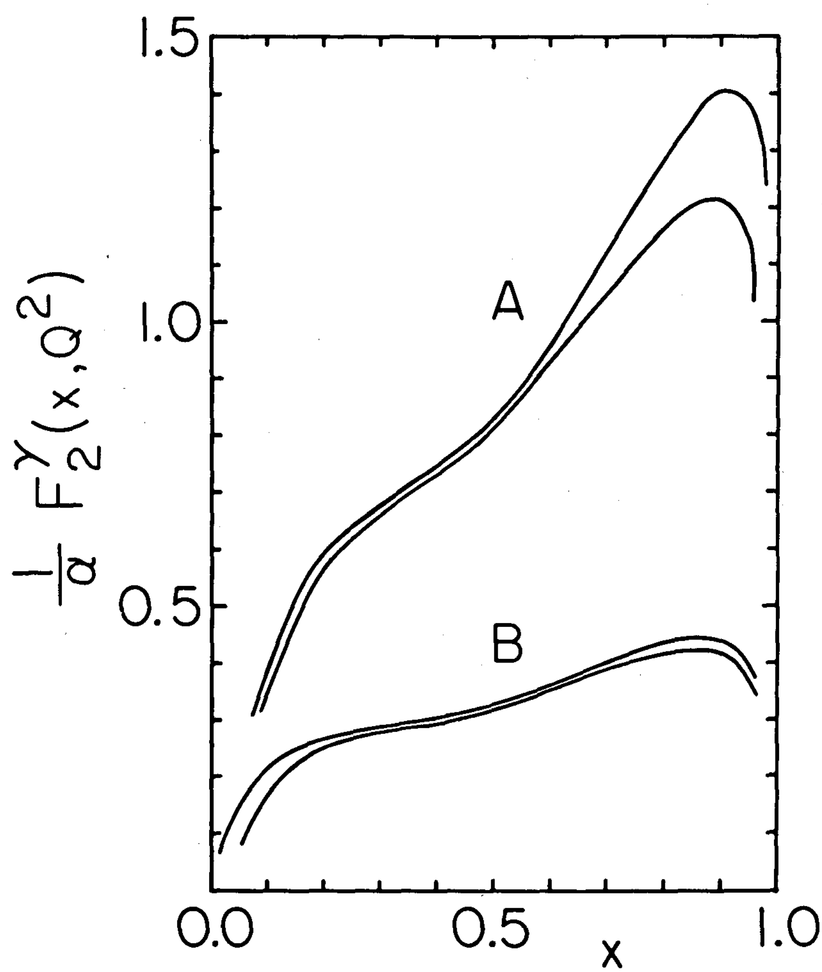

Fig. 5. The result of Uematsu and Walsh for A) Q^2=2000 and B) Q^2=20 GeV^2. The upper(lower) curve in each case corresponds to Λ=100(500) MeV.

summarizes in a simple way our nonperturbative ignorance. An example of their results is shown in fig. 6, where the main point to observe is that although we have solved the negativity problem, we unavoidably induce a significant uncertainty in the prediction. In this case the uncertainty is still concentrated in $x < 0.1$ where we had the problem in the first place. Optimistically then we may hope that our predictions are still reliable for $x > 0.1$, but this point requires still more discussion and study before a definite conclusion can be reached.

The extension of this discussion to yet higher orders has been initiated by Rossi[23)], who has shown that the pointlike piece has the following structure:

$$\frac{1}{\alpha_s(Q^2)}\,\frac{a_n}{d_n+1} + \frac{b_n}{d_n} + \alpha_s(Q^2)\frac{c_n}{d_n-1} + \alpha_s^2(Q^2)\frac{e_n}{d_n-2} + \dots \quad .$$

In this expression the denominators d_n always involve only the the one-loop anomalous dimensions, while b_n, c_n, e_n etc. involve two-loop, three-loop, and four-loop etc. calculations. At present c_n, e_n etc. are not available. The following table locates the positions of the poles in n.

d_n	$n_0^{(-)}$	$n_0^{(NS)}$	$n_0^{(+)}$
-1	1.5926	.3099	-
0	2.0000	1.000	-
1	5.326	5.250	2.386
2	26.52	26.58	4.402

We see that at higher orders the singularities quickly move to large n_0, implying stronger and stronger singularities x^{1-n_0}. Whether or not these singularities contaminate the large-x region is a question of the size of the calculable, but as yet uncalculated, residues.

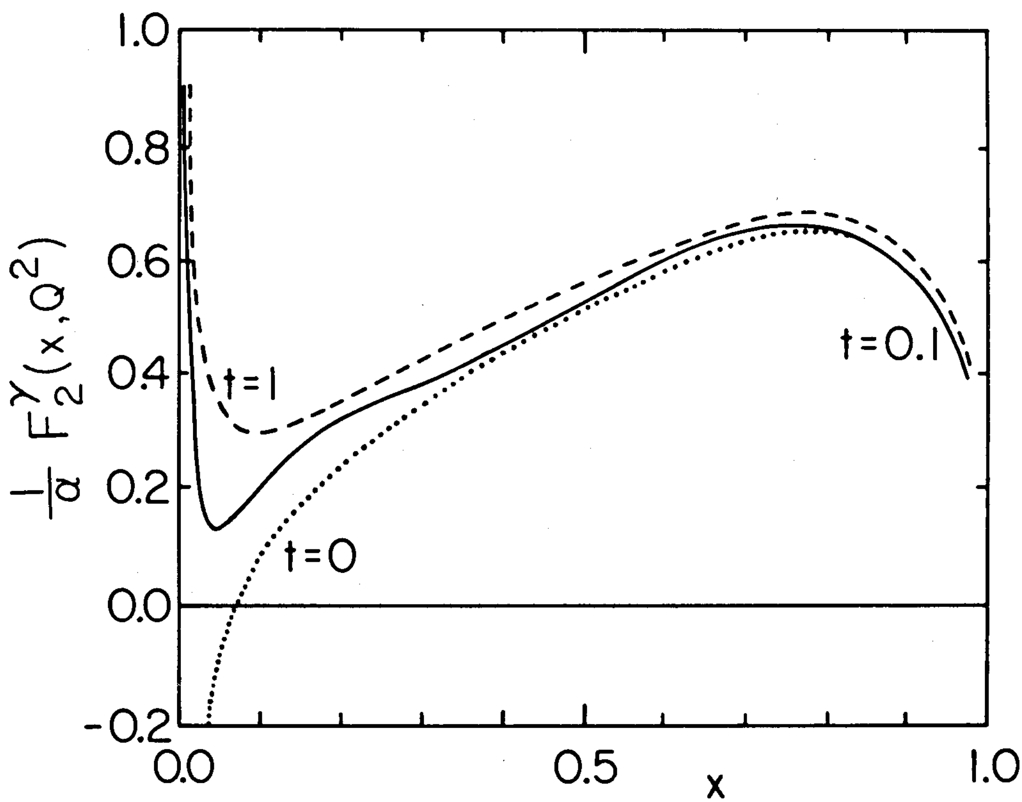

Fig. 6. The reult of Antoniadis and Grunberg for the unregularized case(t=0, dotted line) and two values(t=0.1,solid line and t=1.0,dashed line) of the regularization parameter.

B. x~1

For large enough n and fixed Q^2, the explicit calculation shows that b_n^{NS} and hence $M_n^{\gamma}(Q^2)$ becomes negative. Hence at large x, $F_2^{\gamma}(x,Q^2)$ is negative, which is a nonsensical prediction for a cross section.

The reason for this bad behavior is illustrated in fig. 7 for the Uematsu-Walsh case, although the general trend follows for the $p^2=\Lambda^2$ case as well. The valence piece has the general form from Eq.(3)

$$M_n^V \sim \langle e^4 \rangle \quad \frac{a_n}{\alpha_s} + b_n' + B_n^{\gamma}$$

where

$$\frac{a_n}{\alpha_s} + b_n' \sim C_n^{NS}\ X_n^{NS}$$

and

$$B_n^{\gamma} \sim C_n(1,\bar{g}^2,\alpha).$$

In x space $a_n/\ _s + b_n'$ becomes curve A of fig. 7 while B_n^{γ} becomes (see curve B)

$$B^{\gamma}(x) \sim x\left[(1-2x+2x^2)\ \ln(1/x^2)\ -2+6x-6x^2\right].$$

This is just the box graph evaluated for $m^2=0$, $p^2\neq 0$ and $Q^2=p^2$, and it clearly is negative for x>0.6 even though it is still a cross section[24]. The negativity in this case results simply from dropping $O(p^2/Q^2)$ and $O(m^2/Q^2)$ factors which would insure the correct threshold behavior. In an OPE analysis these factors would arise from higher twist contributions. We see that the theoretical predictions at large x are probably very complicated and potentially unreliable, at least at the leading twist two level.

C. Heavy Quarks

The situation regarding heavy quark contributions to F_2^{γ} is really no different in principle than the treatment in the usual nucleon target deep inelastic scattering[25]. The quantitative effect, however, is much larger for F_2^{γ}, especially for c and eventually t quarks. Within the OPE heavy quarks have been treated by Hill and Ross[26], but their results have not been used phenomenologically so far, and apparently the subject is plagued by subtleties[27]. The more mundane approach of simply using the box graph is being improved by including first order

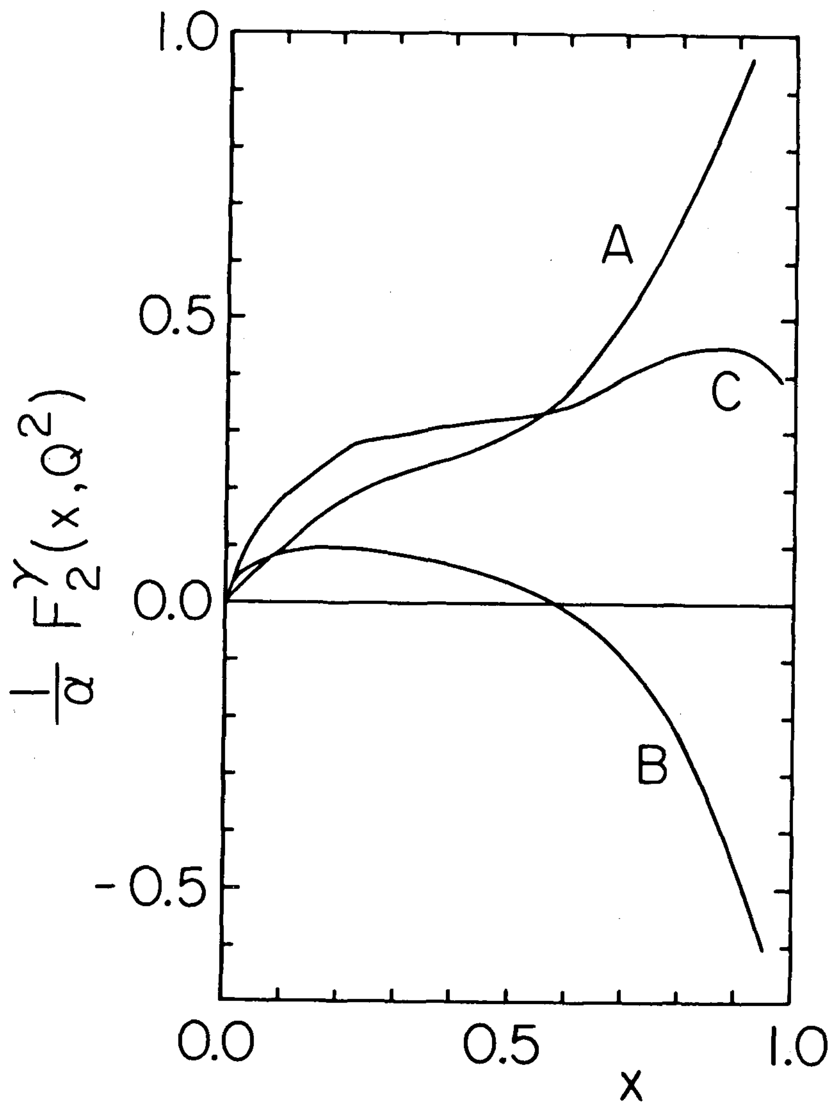

Fig. 7. The breakdown of the valence piece into its Q^2 dependent (curve A) and Q^2 independent (curve B) parts. Curve C is the total.

QCD effects and should be of some use at least not too far above threshold[28].

Hopefully this area will soon receive more theoretical attention.

V. Summary

The hope that experimental measurements of F_2^γ will provide exceptionally "clean" tests of QCD is, of course, still very much alive and worthy of continuing strong experimental and theoretical effort. It is surely no accident and indeed very encouraging that all the experimental results are in beautiful qualitative agreement with theoretical expectations. There is good reason, therefore, to believe that we are observing the predicted dominance of the pointlike coupling of the target photon to its quark constituents.

As a method to measure Λ , we are apparently not so well off as previously thought. We are probably reduced, for finite values of Q^2, to emulating the program of deep inelastic nucleon scattering, i.e. measuring Λ via the Q^2 evolution of the structure functions. This would require using as input to the theory boundary conditions taken from experiment or, more dangerously, some model such as VMD. The feasibility of such a program for realistic experimental situations is not entirely obvious and is under investigation.[29]

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DISCUSSION

Q. J. H. Field - Paris

Is not your final conclusion that F_2^γ is not so useful for measuring Λ unduly pessimistic ?

After all, for the case of the nucleon structure function perturbative QCD is able to calculate only the Q^2 dependence of relatively small corrections to a structure function which is a priori uncalculable. On the other hand, for the photon structure function there are uncalculable (but probably small) corrections to a structure function which *is* calculable both in absolute shape and in terms of Q^2 evolution by perturbative QCD.

A. D.W. Duke

For truly asymptotic values of Q^2, your reasoning is correct. For $Q^2 < 100$ GeV^2 or so, where the transition from Eq.(4) to Eq.(5) of the text is not so harmless, the shape of F_2^γ looks about right but the absolute normalization of the theoretical prediction has been lost, not to be regained until Q^2 is very large.

Q. S. Brodsky - SLAC

Consider fixed quark mass m with $x, Q^2 \neq 0$ and let $\Lambda \to 0$. Does this give the box answer ? Assuming this is the case, perhaps the charm component of the photon structure function can be used to check QCD, although not to measure Λ.

A. D. W. Duke

I think that the answer to the question is yes, and I think that developments along this line will be very important to further test QCD.

Q. Ch. Berger - Aachen

Do you encourage us to go through the pain of measuring the Q^2 dependence of F_2^γ ?

A. D. W. Duke

Even under the not so optimistic circumstances that I have discussed, the photon structure function still enjoys a very favored position in all of perturbative QCD. It is, as far as we know, one of the quantities least sensitive to hadronization and other such nonperturbative and/or uncalculable effects. So every effort should be made to measure F_2^γ as well as possible.

C. W. A. Bardeen - Fermilab

I would like to comment on the theoretical status of the photon structure function. Duke has emphasized the existence of singularities in the pointlike component which must be cancelled by similar singularities in the hadronic component. These singularities were responsible for the large higher order corrections found at small x by Bardeen and Buras. The cancellation of these large effects by the hadronic component should actually improve the predictive power of perturbative QCD as the higher order terms in α_s are now properly suppressed relative to the leading orders. Similar singularities encountered in the pointlike component in yet higher orders should also be suppressed by analogous cancellations. During the next few years the remaining ambiguities in the hadronic component may be resolved through real solutions to QCD (lattice calculations, etc.) which go beyond perturbation theory.

Finally, I think the computation of the charm quark contribution including perturbative corrections should be reliable at low Q^2 and can be matched to the asymptotic forms using reasonable physical analysis at high Q^2.

The measurement of the photon structure function has already shown remarkable agreement in shape and in magnitude with the fundamental predictions of perturbative QCD. In the future we can expect great improvement in both our experimental and our theoretical understanding of these processes.