DEUTSCHES ELEKTRONEN-SYNCHROTRON DESY

DESY 80/52
June 1980

PHOTON-PHOTON PHYSICS IN THE DEEP INELASTIC REGION

by

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Photon-Photon Physics in the Deep Inelastic Region

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Hamburg

Workshop on $\gamma\gamma$ Interactions
Amiens, France
April 8-12, 1980

Abstract

The topics discussed are

1. The real photon structure functions
2. Theoretical issues in $\gamma\gamma$ reactions with one or both photons off mass shell
3. Experimental questions

At this meeting we heard a report by Berger on the observation of inelastic electron photon scattering at PETRA.[1] In the near future we can expect measurements of F_i^γ at Q^2 around 2-5 GeV^2. For a given detector, the accessible Q^2 range increases roughly as E_B (the beam energy). So in the distant future, the F_i^γ could be measured at LEP for $Q^2 \approx$ 15-20 GeV^2. (This time scale is not unnatural; changes in this field occur typically on a scale $\sim$ 5 years.) So it is a good time to recapitulate. This is a brief discussion of what we know theoretically about real photon structure functions and why it is interesting to measure them.

The photon structure functions are a very clean laboratory for theorists. QCD is an elegant theory, but not a trivial one. There is a lot which is only partially understood. Deep inelastic reactions are at present certainly under the best theoretical control. We need experiments on something besides the proton and neutron. Now the photon target is slowly becoming available.

I. REAL PHOTON STRUCTURE FUNCTIONS

The cross section for scattering of an e^+ or e^- on a real photon target in $e^+e^- \longrightarrow e^+e^-$ + hadrons (Fig. 1)[2]

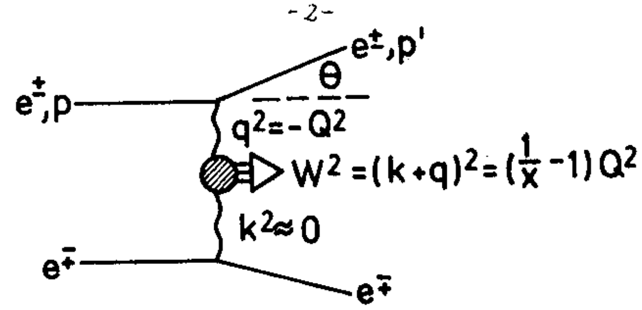

Fig. 1

is

$$\frac{d\sigma}{dx\,dy\,d\phi/2\pi} = \frac{4\pi\alpha^2\, p\cdot k}{Q^4}\left[1 + (1-y)^2\right] \cdot \left\{2x F_T^\gamma + \epsilon(y) F_L^\gamma + \epsilon(y)\,\epsilon\!\left(\frac{E_\gamma}{E_B}\right) F_X^\gamma \cos 2\phi\right\} \qquad (1)$$

where x and y are neutrinolike variables

$$x = Q^2/2p\cdot k\,, \quad y = q\cdot k/p\cdot k$$

and

$$\epsilon(y) = \frac{2(1-y)}{1+(1-y)^2}$$

Defined similarly to the conventional structure functions,

$$F_1^\gamma(x,Q^2) = F_T^\gamma(x,Q^2)$$
$$F_2^\gamma = 2x\,F_T^\gamma + F_L^\gamma \qquad (2)$$
$$F_3^\gamma = F_X^\gamma$$

F_3 is actually the structure function for a transversely polarized photon (analogous to the cross section difference $\sigma_\parallel - \sigma_\perp$ for parallel and perpendicular $\gamma(k)$, $\gamma(q)$ transverse polarizations). It is only measurable if the small angle electron in Fig. 1 is tagged at some non-zero angle. ϕ is then the angle between the planes of the scattered large and small angle $e^\pm$. Actually, for non-zero tagging angle, $k^2 \neq 0$.

So we are ignoring other small contributions when referring to only 3 structure functions in (1).[3] (F_1 and F_2 are measurable without tagging - and therefore for tiny $k^2 \sim O(-m_e^2)$ - if W^2 can be reconstructed by a central detector like PLUTO.)

If the photon were merely a hadron $\sim e^2/f_\rho^2 \sim 1/300$ of the time, as once seemed to be the case, its structure functions would be (at some large Q_o^2)

$$F_T^\gamma(x,Q_o^2)\Big|_{HAD} \approx \left(\frac{e}{f_\rho}\right)^2 F_T^{\rho^0}(x,Q_o^2) \approx \frac{1}{2x}\left(\frac{e}{f_\rho}\right)^2 \frac{1}{4}(1-x)$$

$$F_L^\gamma(x,Q_o^2)\Big|_{HAD} \approx O\left(\frac{\langle p_\perp^2\rangle}{Q_o^2}\right) + O\left(\alpha_s(Q_o^2)\right) \qquad (3)$$

$$F_X^\gamma(x,Q_o^2)\Big|_{HAD} \approx O\left(\frac{\langle p_\perp^2\rangle}{Q_o^2}\right) + O\left(\alpha_s(Q_o^2)\right)$$

(The expression for $F_T^{\rho^0} = F_T^{\pi^0}$ from the quark model and π structure function"data" in ref. (4).) The two terms in $F_{L,X}^\gamma$ are the parton model expressions $O(\langle p_\perp^2\rangle/Q_o^2)$ and the QCD radiative corrections, which are first order in $\alpha_s(Q^2)$. Beyond Q_o^2, F_T^ρ evolves slowly with Q^2, decreasing as an inverse power of $\ln Q^2$ at any fixed x. In any case, $F_T^{\rho^0}$ vanishes as $x \to 1$.

If (3) were all there was to photon structure, it would still be interesting. But (3) is not all there is. The real photon can disassociate at a point into a $q\bar{q}$ pair, one of which is hit by the virtual photon (Fig. 2). The contribution of the box diagram to this pointlike piece is[5]

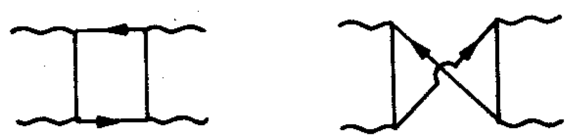

Fig. 2

$$F_T^\gamma(x,Q^2)\Big|_{BOX} = \frac{\alpha}{2\pi}\sum e_i^4\left(x^2+(1-x)^2\right)\ln W^2/\Lambda^2$$

$$\longrightarrow \frac{\alpha}{2\pi}\sum e_i^4\left(x^2+(1-x)^2\right)\ln Q^2/\Lambda^2$$

$$F_L^\gamma(x,Q^2)\Big|_{BOX} = \frac{4\alpha}{\pi}\sum e_i^4 x^2(1-x) \qquad (4)$$

$$F_X^\gamma(x,Q^2)\Big|_{BOX} = -\frac{\alpha}{\pi}\sum e_i^4 x^3$$

This (parton) contribution already shows the important physics: it dominates as $x \to 1$ and $F_2 \propto \ln Q^2$. (Also, F_L^γ and F_X^γ scale.)[5]

F_2^γ contains two terms. The one proportional to $(1-x)^2$ is a $\gamma(q)+\gamma(k) \to \gamma(q)+\gamma(k)$ helicity amplitude where the total spin projection along the collision axis is $J_Z = \pm 2$ (helicities $+- \to +-$ or $-+ \to -+$). The term proportional to x^2 has $J_Z = 0$ along the collision axis ($++ \to ++$ or $-- \to --$). In the limit $Q^2 = -q^2 \to 0$ $x \to 0$ and $J_Z = \pm 2$ dominates. For $x \to 1$ (i.e. fixed W^2 as $Q^2 \to \infty$), $J_Z = 0$ dominates. The argument of the log is $W^2 = (\frac{1}{x} - 1)Q^2$ because it arises from a phase space integral. This is bounded by the CM energy W.

<u>An Aside on Resonances</u>

It's instructive to look at the contribution of a single resonance to F_i^γ. The π^0 final state is a good example, because all others turn out to behave similarly.[6] F_L vanishes and

$$F_{\binom{T}{X}} = \binom{+}{-}\frac{1}{4}\,\delta(W^2 - m_\pi^2)\,|T^{++}|^2 \qquad (5)$$

where

$$T^{++} = \frac{1}{2}Q^2 F_\pi(q^2, k^2=0)$$

(The contribution of a normal or abnormal parity resonance to F_x is + or -). It turns out that in the Bjorken-Johnson-Low limit

$$\lim_{-q^2 = -k^2 \to +\infty} T^{++} = \frac{-\sqrt{2}f_\pi}{3} \qquad (6)$$

Moreover, one can write an integral representation for $F_\gamma(q^2,k^2)$ in terms of a spectral weight. If the weight is concentrated around equal q^2 and k^2, T^{++} is smooth as $k^2/q^2 \to 0$ and [6]

$$T^{++}(q^2,0) \approx -\frac{\sqrt{2}f_\pi}{3} \tag{7}$$

This scaling behavior is independent of QCD radiative corrections (there are none in the BJL limit). The behavior (7) is what one would expect if, from simple duality ideas, the resonances modulated the $x \to 1$ limit of F_T and F_X. We expect that both approach finite constants at fixed W^2 and $Q^2 \gg W^2$. Incidentally, this makes it clear that we should not expect $F_T \propto \ln Q^2$ independent of x as $x \to 1$. $F_T \propto \ln W^2$ is more reasonable.

The helicity structure of the resonance contributions is interesting. At $Q^2 = 0$ the $f^0(1250)$ is produced in a $J_Z = \pm 2$ state and with θ_π measured relative to the collision axis,

$$f^0(1250) \to \pi^+\pi^- \quad : \quad \frac{d\sigma}{d\cos\theta_\pi} \propto \sin^4\theta_\pi \tag{8}$$

However, at <u>large</u> Q^2, the f^0 is produced in a $J_Z = 0$ state[6] and

$$f^0(1250) \to \pi^+\pi^- : \quad \frac{d\sigma}{d\cos\theta_\pi} \propto (2 - 3\sin^2\theta_\pi)^2 \tag{9}$$

Of course, this says nothing about the rapidity of the change from (8) to (9).

<u>QCD Radiative Corrections</u>

The box diagram (or parton model) result for the pointlike term in F_T^γ is changed by QCD radiative corrections.[7] The box result, written as an integral over the quark $p_\perp^2$ is of the form [2]

$$F_T^\gamma \sim \alpha \int \frac{dp_\perp^2}{p_\perp^2} \sim \alpha \ln Q^2 \tag{10}$$

Radiation of N gluons will give a contribution of order (Fig. 3)

$$\alpha \int \frac{dp_\perp^2}{p_\perp^2} \alpha_s \int \frac{dp_\perp^2}{p_\perp^2} \cdots \quad \propto \alpha \ln Q^2 (\alpha_s \ln Q^2)^N \tag{11}$$

(The $p_\perp^2$ integrals have nested limits, but as $Q^2 \to \infty$ this does not matter and (11) results). A sum of terms like (11) resembles a geometrical series which sums up to a constant times the lowest order result $\propto \ln Q^2$. Thus the $\ln Q^2$ behavior survives, but we expect that the x-dependence will differ from (4). It will be softer because the quarks loose momentum by gluon radiation.

The QCD corrections to F_L differ. To lowest order in α_s, the integrand to F_L^γ

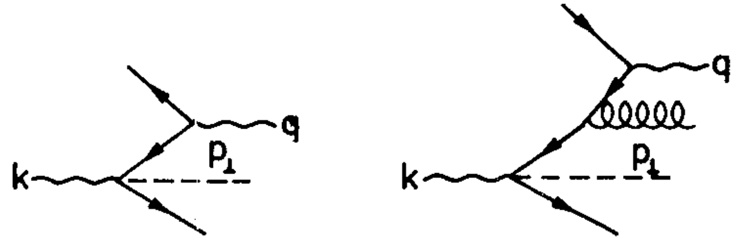

Fig. 3

vanishes at $p_\perp = 0$ for massless quarks. It can be written as $p_\perp^2/Q^2 + O(\alpha_s(Q^2)) +$ higher order terms,

$$F_L^\gamma \sim \alpha \int \frac{dp_\perp^2}{p_\perp^2} \frac{p_\perp^2}{Q^2} + \alpha \int \frac{dp_\perp^2}{p_\perp^2} \alpha_s(Q^2) + \cdots$$
$$\sim \alpha \cdot \text{const} + O(\alpha_s \ln Q^2) + O(\alpha_s^2 \ln Q^2) \tag{12}$$

The first two terms give constants as $Q^2 \to \infty$; all others vanish as powers of $1/\ln Q^2$. The numerical QCD corrections to the first (parton) term turn out to be very small. It is not likely that they will ever be measured (if QCD is correct).

Perhaps startlingly, there are no O(1) QCD corrections to F_X.[8] QCD effects are of order $\alpha_s(Q^2)$, and vanish as $Q^2 \to \infty$. This can be seen by a simple argument. F_X is not actually a cross section. It is the difference of cross sections for parallel and perpendicular polarized $\gamma(k), \gamma(q)$. By symmetry, the probability to find a q in $\gamma(k)$ must vanish at $p_\perp^2 = 0$ – no matter how many gluons have been radiated, provided their momenta and angles are integrated over. This is because $F_X^\gamma \propto \sigma_\parallel - \sigma_\perp$. The angular factor this introduces is $\sin^2\theta_q \propto p_\perp^2/Q^2$. This means that F_X^γ takes the form

$$F_x^\gamma \sim \alpha \int \frac{dp_\perp^2}{p_\perp^2} \frac{p_\perp^2}{Q^2} + \alpha \int \frac{dp_\perp^2}{p_\perp^2} \alpha_s(Q^2) \frac{p_\perp^2}{Q^2} + \cdots \quad (13)$$

$$\sim \alpha \cdot \mathrm{const} + O(\alpha_s(Q^2))$$

Thus the QCD corrections to the box diagram result vanish as $Q \rightarrow$.

Fig. 4 shows the behavior of F_2^γ

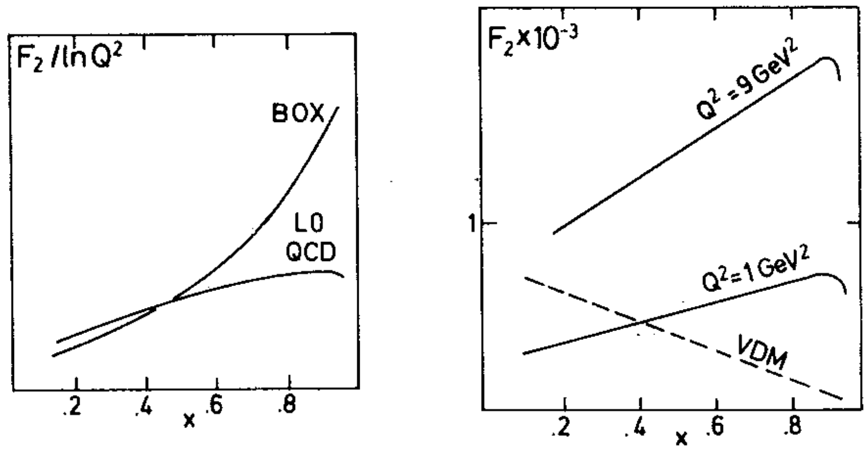

Fig. 4

II. ISSUES WITH ONE OR TWO PHOTONS OFF MASS-SHELL

1) F_T^γ, F_L^γ and F_x^γ probe the $p_\perp$ structure of QCD corrections in different ways. If it is at all possible, all three should be measured.

2) In inelastic electron-photon-scattering there is no target mass. So there are no target mass corrections as in eN scattering. For the pointlike term there are also no corrections from the non-zero radius of the target. We also expect $F_2 \propto \ln Q^2$, rising with Q^2, unlike ℓN scattering, where it decreases. It should prove difficult to explain the Q^2 dependence of F_2 solely as a hadronic higher-twist effect, as appears possible for lepton-nucleon scattering.

3) Higher order QCD corrections to F_2 have been calculated.[9] These can be important

at very large x. Much of the correction may be kinematic in nature. The argument of the log in F_2^γ is not Q^2, but W^2. For $x > 1/2$, $W^2 < Q^2$ and the kinematic limits in $\gamma\gamma \rightarrow q\bar{q} + q\bar{q}G + \ldots$ are set by W^2, not Q^2. (For example propagators can go off shell by an amount $W^2 < Q^2$). This suppresses F_2 at $x > 1/2$, compared to the leading $\ln Q^2$ result. At small x, where $W^2 \gg Q^2$, the hadronic structure of γ dominates and no clear statement can be made. This kinematic effect is clearly nonleading, as $\ln W^2 = \ln Q^2 + \ln(1/x - 1)$.

4) QCD predicts a simple relation between $e\,\gamma(k) \rightarrow e' X$ and $e^+e^- \rightarrow \gamma(q) \rightarrow \gamma(k) + X$. In the parton model this is trivial, and it is straightforward even including QCD corrections. Moreover, $e^+e^- \rightarrow \gamma(k) + X$ should be measurable at very large $Q^2 \sim 100\ \mathrm{GeV}^2 - 1000\ \mathrm{GeV}^2$ at CESR and PETRA or PEP.[10]

Futurism

F_T is well-defined at LEP, where one can define an off-shell extension

$$F_T^\gamma(Q^2, |k^2|) = \frac{1}{2}\left[W_{++\rightarrow++} + W_{+-\rightarrow+-}\right] \quad (14)$$

In the limit $|k^2|$, Q^2, $W^2 \rightarrow \infty$ with fixed $x = 1 + W^2/Q^2$ and $Q^2/|k^2|$ very large there are no QCD radiative corrections and

$$F_T^\gamma(Q^2, |k^2|) \xrightarrow[|k^2| \ll Q^2]{} \frac{\alpha}{2\pi} \sum e_i^4 \left(x^2 + (1-x)^2\right) \ln \frac{Q^2}{|k^2|} \quad (15)$$

the box diagram (or parton) result. By increasing $|k^2|$ from zero one "turns off" the QCD radiative corrections. Presumably the scale for this is set by the QCD dimensional parameter Λ.

It will clearly be interesting to measure (14) at LEP to see how fast the QCD radiative corrections disappear. (As yet there are no detailed theoretical calculations of this effect known to me.[11])

Another thing which may barely be measurable at LEP is the large $|k^2|$, $|q^2|$ behavior for fixed W^2. QCD radiative corrections do survive in this limit. The box result is modified.[12]

III. EXPERIMENT

There is clearly a background to eγ scattering in Fig. 1. It is the inelastic compton process of Fig. 5.[13]

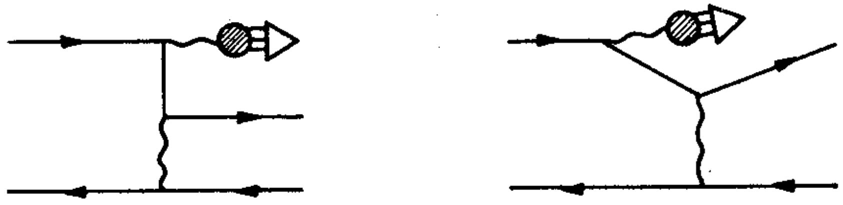

Fig. 5

For small angle scattering, where an e^- is tagged at some Q^2 and not too large angle to the e^- beam, this background is harmless. The reason is that the propagators in Fig. 5 are both far off-shell. This is not so if one does not distinguish e^+ and e^- charges in the final state. Then one might see an e^+ at small or moderate angle to the e^- beam (and the reverse). This is not obviously harmless, because the first graph in Fig. 5 has a u-channel pole (backward scattering). However, the outcoming lepton has small laboratory energy. An energy cut on the outgoing lepton will reduce this problem. Fig. 6 shows the cross section as a function of E' and θ of the scattered electron (or, as a dashed line, electron plus backscattered e^+ if charges are not distinguished). With a cut on the scattered e^- energy of $E' \gtrsim 2$ GeV or $\gtrsim 12$ GeV the inelastic Compton background is unimportant below $\approx 20^o$ scattering angle. This incorporates the whole angular range of the PLUTO "large angle" tagger.[2] The situation is essentially the same at LEP.[2]

As a final remark, it is worth noting that F_L and F_X may be measurable. F_L is larger in eγ scattering than in ℓN.[2] The relevant distributions are of the form

$$[1+(1-y)^2]F_2^\gamma - y^2 F_L^\gamma \equiv A - B$$

$$[1+(1-y)^2]F_2^\gamma - y^2 F_L^\gamma + 2(1-y)F_X^\gamma \cos 2\phi \qquad (16)$$

$$\propto 1 - a\cos 2\phi$$

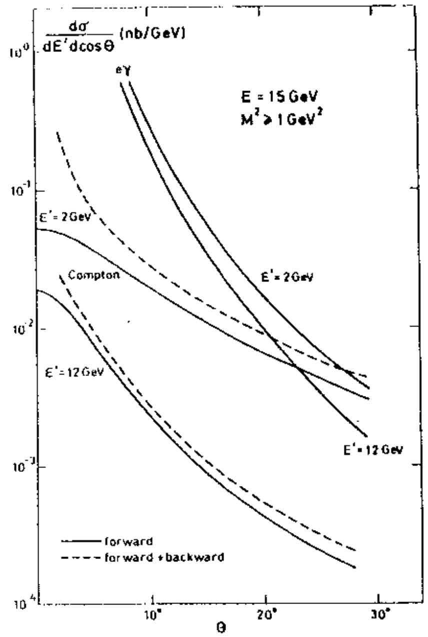

Fig. 6

Numerically, B/A and a are large enough to offer some hope that a high statistics experiment can extract F_L^γ and F_3^γ.

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