Quarks and gluons in nuclei

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Abstract

We give an overview of current results and problems in deep inelastic scattering of leptons on nucleons and nuclei

I. Deep Inelastic Lepton-Nucleus Scattering

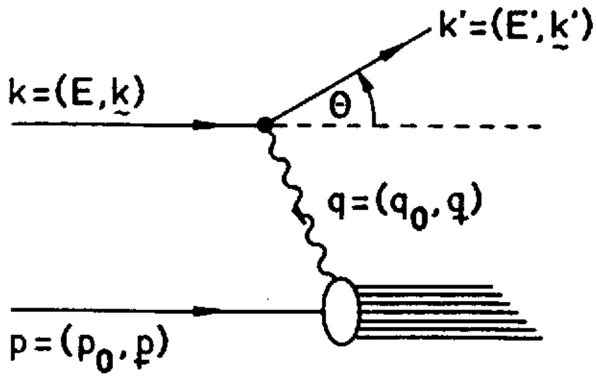
PARTON MODEL OF THE NUCLEON

The most natural way of studying the nucleus in detail is to look at it with a high resolution microscope. Nowadays high energy electron, muon or neutrino beams are the best instruments to investigate hadronic structure, at short distances. The resolution Δ of these microscopes is related to the maximal momentum transfer Q

$$\Delta \sim \hbar/Q \tag{1.1}$$

which is a function of scattering angle and incoming energy. In the post-war history of nuclear and particle physics we saw lepton scattering experiments with increasing resolving power. In the 1950s Hofstadter (1963) and his collaborators studied the sizes of nuclei with elastic electron scattering. In the 60s SLAC's (Stanford Linear Accelerator) new GeV accelerator was used to study the charge and current distribution of the proton. In Cornell and DESY many interesting experiments were done on nuclei, especially shadowing (Grammer, jr. and Sullivan, 1978) due to the hadronic component of the photon was measured at GeV energies. Only in the 1970s (Bloom, 1975; Friedman and Kendal, 1972), with a beam energy of 20 GeV, were the hard constituents of hadronic matter - the quarks - discovered in deep inelastic electron-proton scattering. Very recently at CERN (Drees and Montgomery, 1983) muons have been used to map out the motion and distribution of partons (quarks and gluons) in hadrons. Muons have the advantage that radiative QCD corrections are less important. Deep inelastic lepton scattering experiments were of fundamental significance for the establishment of Quantum Chromodynamics (QCD) as the theory of strong interactions (Wilczek , 1982). A change of parton structure function with photon resolution ΔQ^2 can be attributed to a weak, i.e. asymptotically vanishing QCD-interaction between quarks and gluons. It is this success of parton models which has triggered also the development of bag models (Chodos and others, 1974; Thomas, 1982) for the static structure of baryons. In these bag models free quarks are imprisoned in a cavity by the QCD-vacuum fluctuations. How the confinement of quarks (Hooft, 1980) really works is an intensive topic of actual research. Lattice gauge theory (Kogut, 1982) seems to be the most promising method to demonstrate confinement and calculate the mass spectrum of hadrons. However, it is still not clear how to combine the picture of the proton obtained in deep inelastic scattering with the static description of the proton.

The electromagnetic interaction of a lepton with a hadron is mediated by the ex-



$$Q^2=-q^2=(k-k')^2$$
$$\nu=E-E'=p\cdot q/m$$
$$x=Q^2/2m\nu$$

Fig. 1.1. Basic kinematic variables for deep inelastic lepton scattering.

change of a virtual photon with four-momentum q. The basic variables of the photon are its four-momentum square (e.g. Fig. 1.1)

$$Q^2 = -q^2 = -(k - k')^2 = 4\ EE'\ \sin^2\theta/2 > 0 \tag{1.2}$$

and its energy

$$\nu = E - E' \tag{1.3}$$

where E, E' (k, k') are the laboratory energies (four-momenta) of the incident and scattered lepton. The photon can have transverse (helicity ± 1) or longitudinal (h = 0) polarization. The momentum and mass of the hadron are p and m. One can re-express the lepton-energy loss in the laboratory with Lorentz-four vectors as

$$\nu = p\cdot q/m. \tag{1.4}$$

Instead of ν it is also possible to choose W^2, the mass squared of the hadronic system after scattering, as second variable:

$$W^2 = (p + q)^2 = m^2 + 2m\nu(1 - \frac{Q^2}{2m\nu}) \tag{1.5}$$

or the Bjorken variable $x = Q^2/2m\nu$ such that

$$W^2 = m^2 + Q^2(\frac{1}{x} - 1). \tag{1.6}$$

For elastic scattering $W^2 = m^2$, i.e. $x = 1$. Because of eq. (1.2) $0 \leq x \leq 1$.

In general the inelastic cross-section of a lepton (e,μ) from an unpolarized target can be reduced to two structure functions due to the two polarizations of the exchanged virtual photon. The polarization vector of the photon is $\varepsilon(\lambda)$ with helicity $\lambda = \pm 1$ for transverse photons and helicity $\lambda = 0$ for longitudinal scalar photons+ is given as

$$\varepsilon(\lambda = \pm 1) = \mp \frac{1}{\sqrt{2}} (0,\ 1,\ \pm i,\ 0),$$
$$\varepsilon(\lambda = 0) = \frac{1}{\sqrt{Q^2}} (\sqrt{Q^2 + \nu^2},\ 0,\ 0,\ \nu), \tag{1.7}$$

where we have chosen the three-momentum of the photon in z-direction $q = (\nu,\ 0,\ 0,\ |q|)$. The above vectors are normalized to $\varepsilon_\mu \varepsilon^\mu = \varepsilon \cdot \varepsilon = 1$ and satisfy gauge invariance which demands $\varepsilon \cdot q = 0$. The scattering amplitude for inelastic scattering of a target T leading to a final state X is (Bjorken and Drell, 1964; Close, 1979)

$$S_{fi} = i(2\pi)^4 \delta(k + p - k' - p_X) \bar{u}(k') \gamma^\mu u(k) \frac{e^2}{q^2} < p_X | J_\mu | p >. \tag{1.8}$$

J_μ is the hadronic component of the electromagnetic current. For unpolarized lepton and target the inclusive cross-section $\ell + T \rightarrow \ell' + X$ becomes for relativistic lepton energies $E >> \mu$ (μ = mass of the lepton, M = mass of the target nucleus)

$$d\sigma = \frac{\mu^2}{EE'} \frac{d^3k'}{(2\pi)^3} \sum_X (2\pi)^4 \delta^4(k + p - k' - p_X) (\frac{e^2}{q^2})^2$$
$$\frac{1}{2} \mathrm{Tr}(\gamma^\mu \frac{\not k}{2\mu} \gamma^\nu \frac{\not k'}{2\mu}) \sum_{Pol} < p | J_\mu(0) | p_X >< p_X | J_\nu(0) | p >. \tag{1.9}$$

This formula is derived in Bjorken-Drell (Bjorken and Drell, 1964) for a current J_μ given by a massive spin-1/2 particle. The $\sum_{Pol}$ includes the averaging over initial polarizations. The trace over the lepton spin gives

$$\frac{1}{2} \mathrm{Tr}(\gamma^\mu \frac{\not k}{2\mu} \gamma^\nu \frac{\not k'}{2\mu}) = \frac{1}{2\mu^2} (k^\mu k'^\nu + k^\nu k'^\mu - g^{\mu\nu} kk'). \tag{1.10}$$

All the physics of the hadron is contained in the current-current correlation function

$$W_{\mu\nu} = \sum_X (2\pi)^4 \delta^4(p_X - p - q) \sum_{Pol} < p | J_\mu(0) | p_X >< p_X | J_\nu(0) | p >$$
$$= \sum_{Pol} \int d^4x\, e^{iqx} < p | J_\mu(x) J_\nu(0) | p > \tag{1.11}$$

using

$$< p | O(x) | p_X > = e^{-i(p_X - p)x} < p | O(0) | p_X >. \tag{1.12}$$

Because of the symmetry under exchange of μ and ν in eq. (1.10) $W_{\mu\nu}$ has to be symmetric, too. Gauge invariance demands $q^\mu W_{\mu\nu} = 0$; consequently $W_{\mu\nu}$ depends only on two unknown functions $W_1(\nu, q^2)$ and $W_2(\nu, q^2)$ (M = mass of the target nucleus)

+Note: Four vectors are written as (a_0, a_1, a_2, a_3) if not indicated otherwise.

$$W_{\mu\nu} = W_1(\nu,q^2)\ (\frac{q_\mu q_\nu}{q^2} - g_{\mu\nu}) + W_2(\nu,q^2)\ \frac{1}{M^2}\ (p_\mu - \frac{pq}{q^2}\ q_\mu)(p_\nu - \frac{pq}{q^2}\ q_\nu). \tag{1.13}$$

The cross-section can then be expressed with $(\frac{d\sigma}{d\Omega})_{Mott}$, the Mott cross-section for the scattering of a relativistic electron in a Coulomb field (Bjorken and Drell, 1964)

$$(\frac{d\sigma}{d\Omega})_{Mott} = \frac{4\alpha^2\ E'^2\ \cos^2\theta/2}{Q^4} \tag{1.14}$$

as

$$\frac{d\sigma}{d\Omega dE'} = (\frac{d\sigma}{d\Omega})_{Mott}[W_2 + 2W_1 tg^2\theta/2]. \tag{1.15}$$

In general it is not easy to measure large angle scattering. Therefore W_2 is better known than W_1.

Using the polarization vectors $\varepsilon(\lambda)$ (eq. (1.7)) it is conventional to define longitudinal and transverse response functions S_L and S_T in nonrelativistic nuclear physics (Donnelly and Walecka, 1975)

$$\begin{aligned} S_L(\nu,Q^2) &= \frac{\vec{q}^2}{Q^2}\ \varepsilon_o^{*\mu}\ \varepsilon_o^{\nu}\ W^{\mu\nu} = \frac{\vec{q}^2}{Q^2}(-W_1 + \frac{\vec{q}^2}{Q^2}\ W_2) \\ S_T(\nu,Q^2) &= \sum_{\lambda=\pm 1}\ \varepsilon_\lambda^{*\mu}\ \varepsilon_\lambda^{\nu}\ W^{\mu\nu} = 2W_1. \end{aligned} \tag{1.16}$$

Note the additional factor $\vec{q}^2/Q^2$ in eq. (1.16) for the nonrelativistic definition of $S_L(\nu,Q^2)$. Then the double differential cross-section has the form ($Q^2 = -q^2$)

$$\frac{d\sigma}{d\Omega dE'} = (\frac{d\sigma}{d\Omega})_{Mott}[S_L(\nu,Q^2)(\frac{Q^2}{\vec{q}^2})^2 + S_T(\nu,Q^2)(tg^2\theta/2 + \frac{Q^2}{2\vec{q}^2})]. \tag{1.17}$$

In relativistic deep inelastic scattering the virtual scalar-longitudinal σ_S and transverse σ_T photon cross-sections are defined such that (Close, 1979)

$$\begin{aligned} \sigma_T &= \Gamma W_1 \\ \sigma_L &= \Gamma(-W_1 + \frac{Q^2+\nu^2}{Q^2}\ W_2). \end{aligned} \tag{1.18}$$

Γ contains the photon flux $K \approx \nu + Q^2/2\nu$ and the coupling constant α, i.e. $\Gamma = 4\pi^2\alpha/K$.

The most important consequence of the high energy experiments was the existence of point-like substructure in the proton. These experiments were performed at SLAC around 1967 - 1970. A review of the results is given by Bloom (1975) and by Friedman and Kendal (1972). Theoretically the naive quark-parton model of the nucleon can describe the essential features of deep inelastic lepton-nucleus scattering. Its main assumptions are the following ones:

(i) partons (= quarks and gluons) behave like free particles,

(ii) a fast-moving hadron looks like a jet of partons moving in the same direction as the hadron.

The assumption (ii) is based on the cut-off of transverse momenta [$p_\perp^c \sim 300$ MeV] observed in all hadronic reactions. Hypothesis (i) is related to the asymptotic freedom of Quantum Chromodynamics. Stated differently the interaction time of the photon in deep inelastic scattering is so short that the partons are quasifree. This assumption is similar to the Fermi Gas Model of nucleons in the nucleus. When the energy transfer $\nu \gg \hbar w_{\text{shell model}}$, we can neglect the residual interaction. In the same way we expect when $\nu \gg \hbar w_{\text{quark model}} \approx 600$ MeV that the incoherent summation over final states is valid.

The ideal formulation of the parton model is in a reference system, where the bound state moves very fast with P_∞. Intuitively because of Lorentz time dilatation, the internal motion of the constituents is slowed down. The fast-bound state momentum P_∞ can serve as a reference scale if the transverse momenta of the constituents are limited, then in the limit

$$P_\infty \gg \sqrt{<\vec{p}_\perp^{\,2}> + m^2} \tag{1.19}$$

one can expand the "nasty" relativistic energy operator of each constituent i around the z-component of its momentum

$$p_{z,i} = \eta_i P_\infty \tag{1.20}$$

and obtain

$$H_i = \sqrt{(\eta_i P_\infty)^2 + \vec{p}_{\perp,i}^{\,2} + m^2} \simeq \eta_i P_\infty + \frac{\vec{p}_{\perp,i}^{\,2} + m^2}{2\eta_i P_\infty} . \tag{1.21}$$

This expansion allows to represent the total Hamilton operator as a <u>sum</u> of operators H_i,

$$H = \sum_i H_i = P_\infty + \sum_{i=1}^{A} \frac{\vec{p}_{\perp,i}^{\,2} + m^2}{2\eta_i P_\infty} \tag{1.22}$$

which look very similar to operators in nonrelativistic Schrödinger mechanics. Since we can use the conservation law for the momentum fractions of all constituents, we have

$$\sum_i^A p_{z,i} = \sum_i^A \eta_i P_\infty = P_\infty \quad \text{or} \quad \sum_{i=1}^A \eta_i = 1. \tag{1.23}$$

In this infinite momentum frame the wavefunctions will depend on η and $k_\perp$, as $\psi(\eta, k_\perp)$, and evolve with the Hamilton operator of eq. (1.22). A more general framework than the heuristic discussion given above is the concept of light-cone-coordinates. It amounts to a new choice of coordinates rotated by $\pi/4$ from the regular system and has all the advantages of the infinite momentum frame without sharing its limitations (eq. (1.19)).

We define Light-Cone-Coordinates, a new time τ and a new z-axis ξ and the corresponding momenta

$$\begin{aligned} \tau &= t + z & \quad p^+ &= E + p_z \\ \xi &= t - z & \quad p^- &= E - p_z \\ \vec{x}_\perp &= \vec{x}_\perp & \quad \vec{p}_\perp &= \vec{p}_\perp \end{aligned} \tag{1.24}$$

The scalar product $a \cdot b = 1/2(a^+b^- + a^-b^+) - \vec{a}_\perp \vec{b}_\perp$. The integration volume $\int d^4x = \int dt dz d\vec{x}_\perp = \int 1/2 d\tau d\xi d\vec{x}_\perp$. The coordinates of a particle on mass shell are normally

$$\begin{aligned} &(p_0, \vec{p}) && (p^+, p^-, \vec{p}_\perp) \\ &= (\sqrt{\vec{p}^2 + m^2}, \vec{p}_\perp, p_z) && = (p^+, \frac{\vec{p}_\perp^2 + m^2}{p^+}, \vec{p}_\perp), \end{aligned} \tag{1.25}$$

i.e., in the same way as normally the energy is constrained the p^--component is constrained by the on-shell condition. The free wavefunction propagates in light-cone time τ as

$$\begin{aligned} i \frac{\partial}{\partial \tau} |\psi\rangle &= \frac{i}{2}\left(\frac{\partial}{\partial t} + \frac{\partial}{\partial z}\right)|\psi\rangle = \frac{i}{2}\left(\frac{\partial}{\partial t} + \frac{\partial}{\partial z}\right) e^{-ip_0 t + ip_z z} |p_0, \vec{p}_\perp, p_z\rangle \\ &= \frac{1}{2}(p_0 - p_z)|\psi\rangle = \frac{1}{2} p_- |\psi\rangle, \end{aligned} \tag{1.26}$$

i.e. off the p^--shell. The role of the Hamilton operator is played by the p^--operator

$$p^- = \sum_{i=1}^{A} \frac{\vec{p}_{\perp,i}^{\,2} + m^2}{\eta_i P^+} \tag{1.27}$$

which has the same form as eq. (1.22) when one identifies

$$P_+ = \sqrt{P_\infty^2 + M^2} + P_\infty \approx 2P_\infty$$

for large bound-state z-momentum P_∞. For a longer discussion of the light-cone-formalism we refer to Brodsky (1982), Chemtob (1980), Frankfurt and Strikman (1981) and Kogut and Susskind (1973). In Fig. 1.2 we give the derivation of the Bjorken variable for cartesian and light-cone-variables. We neglect the transverse momentum of the partons. In the Breit-frame the initial parton momentum p_z is reversed after the collision with the photon. Applying the parton model to deep inelastic scattering we obtain the current-current correlation function of the proton (eqs. (1.11) and (1.13)) by an incoherent summation of scatterings of the lepton on each parton with charge e_i. Let $N_i(\eta)$ be the probability to find a quark i with momentum fraction η in the proton. Then the current-current tensor $\langle J_\mu J_\nu \rangle$ for massless spin-1/2 quarks is given by

$$\begin{aligned} w_{\mu\nu}^{1/2} &= \mathrm{Tr}_{s,s'} \langle p|J_\mu\, p'\rangle\langle p'|J_\nu|p\rangle = \frac{1}{2}\mathrm{Tr}(\gamma^\mu \not{p}' \gamma^\nu \not{p}) = \\ &= 2(p_\mu p'_\nu + p_\nu p'_\mu - g_{\mu\nu} pp') \end{aligned} \tag{1.28}$$

with $p_\mu = \eta P_\mu$ and $p'_\mu = p_\mu + q_\mu$.

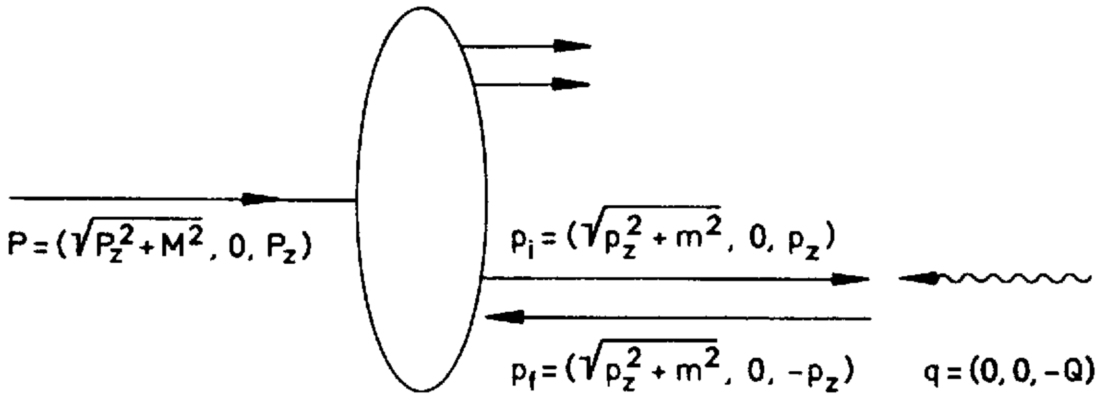
To get $W_{\mu\nu}$ for the proton we add incoherently the quark contributions demanding that the quark in the final state is on the mass shell $m_q^2 \approx 0$:

$$W_{\mu\nu}(\nu, Q^2) = \frac{1}{2m} \sum_i e_i^2 \int \frac{d\eta}{\eta} N_i(\eta)\, w_{\mu\nu}^{1/2}\, \delta((p+q)^2). \tag{1.29}$$

The δ-function yields $\delta(2\eta P \cdot q + q^2) = \delta(2m\nu \cdot \eta - Q^2)$ or for η we get the Bjorken value $\eta = x_B$

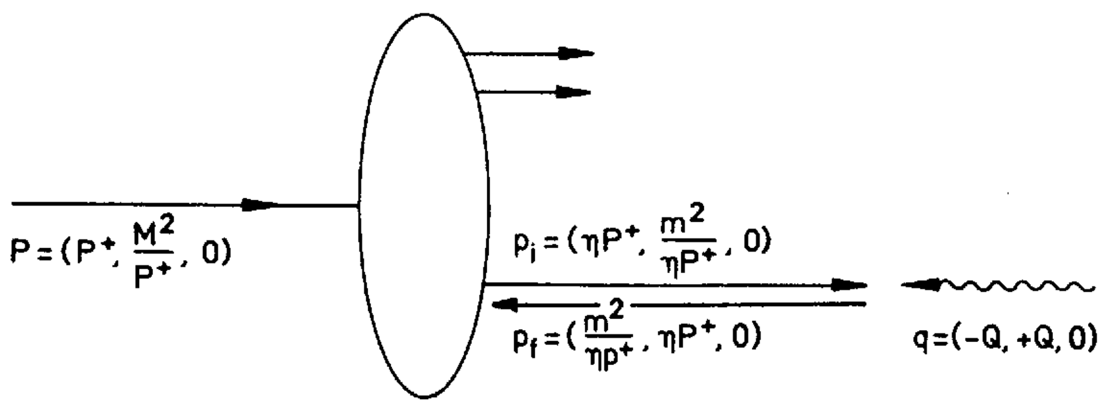
$$x_B = \frac{Q^2}{2m\nu}. \tag{1.30}$$

CARTESIAN COORDINATES: $p = (p_0, p_\perp, p_z)$



$P \cdot q = M\nu = P_z \cdot Q;$ $\quad Q^2 = 4p_z^2;$ $\quad x = Q^2/2M\nu = \frac{p_z}{P_z} \cdot (P_z \gg M)$

LIGHT CONE COORDINATES: $p = (p^+, p^-, p_\perp) = (p_0 + p_z, p_0 - p_z, p_\perp)$



$P \cdot q = M\nu = \frac{1}{2} P^+ Q;$ $\quad Q^2 = \eta^2 P^{+2};$ $\quad x = Q^2/2M\nu = \eta$

Fig. 1.2. Light-Cone-Coordinates versus cartesian coordinates for the absorption of the virtual photon on a constituent in the Breit-frame.

The representation of W_1 and W_2 in terms of the quark distribution function $N_i(\eta)$ is obtained by comparing eq. (1.29) with eq. (1.13) for the proton with mass m and momentum P_μ as target, i.e.

$$W_{\mu\nu} = W_1(\nu,Q^2)\left(\frac{q_\mu q_\nu}{q^2} - g_{\mu\nu}\right) + W_2(\nu,Q^2)\,\frac{1}{m^2}\left(P_\mu - \frac{Pq}{q^2}\,q_\mu\right)\left(P_\nu - \frac{Pq}{q^2}\,q_\nu\right). \tag{1.31}$$

For simplicity we set all terms containing q^μ to zero. We can reconstruct them afterwards from gauge invariance. Then two equations follow from eqs. (1.29) and (1.31)

$$\frac{1}{2m}\sum_i e_i^{\,2} \int \frac{d\eta}{\eta}\, N_i(\eta)\, \frac{1}{2m\nu}\,\delta(\eta - x_B)\{4\eta^2 P_\mu P_\nu - 2g_{\mu\nu}\cdot\eta m\cdot\nu\} =$$
$$= W_2\,\frac{1}{m^2}\,P_\mu P_\nu - g_{\mu\nu}\,W_1. \tag{1.32}$$

Consequently we have the scaling of the inelastic structure functions of Bjorken (1969) for spin-1/2-partons (c.f. Fig. 1.3)

$$\nu W_2(\nu,Q^2) = x_B \sum_i e_i^{\,2}\, N_i(x_B) = F_2(x_B)$$
$$2m\;W_1(\;,Q^2) = \sum_i e_i^{\,2}\, N_i(x_B) = F_1(x_B). \tag{1.33}$$

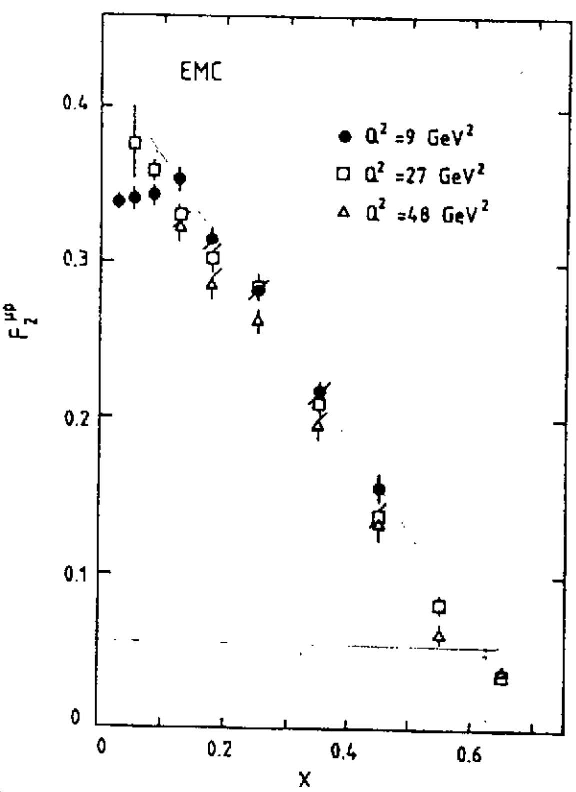


Fig. 1.3. Measurement of the structure function $F_2^{\mu p}(x,Q^2)$ in muon proton scattering (Drees and Montgomery, 1983).

Bjorken scaling states that W_1 and W_2 become in the limit $Q^2 \to \infty$, $\nu \to \infty$ functions of one variable x_B. For spin-0-partons one would find an expression similar to eq. (1.28). Only the current-current tensor $W^{1/2}_{\mu\nu}$ is replaced by

$$W^0_{\mu\nu} = (p_\mu + p'_\mu)(p_\nu + p'_\nu). \tag{1.34}$$

This tensor produces a vanishing W_1 as can be easily seen (eq. (1.32))

$$\begin{aligned} \nu W_2{}^0(\nu, Q^2) &= x_B \,\Sigma\, e_i{}^2\, N_i{}^0(x_B) \\ 2m\, W_1{}^0(\nu, Q^2) &= 0. \end{aligned} \qquad \text{Spin-0-Partons} \tag{1.35}$$

A good way of measuring the spin of the partons is the ratio $R = \sigma_L/\sigma_T$ (eq. (1.18)). For spin-1/2-partons one obtains $R \to 0$, whereas for spin-0-partons the ratio $R \to \infty$. In the scaling limit $\nu, Q^2 \to \infty$ but $Q^2/2m\nu$ fixed R can be expressed by F_1 and F_2

$$R = \frac{\sigma_L}{\sigma_T} = \frac{(1 + \nu^2/Q^2)\, W_2 - W_1}{W_1} \to \frac{F_2 - xF_1}{xF_1}. \tag{1.36}$$

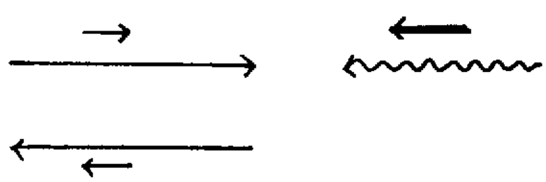
If the fraction of scalar partons weighted with their charges squared is $\gamma(x)$, i.e.

$$\gamma(x) = \frac{\Sigma\, e_i{}^2\, N_i{}^0(x)}{\Sigma\, e_i{}^2\, N_i{}^0(x) + \Sigma\, e_i{}^2\, N_i(x)}, \tag{1.37}$$

one finds for R

$$\begin{aligned} R_{theor} = \frac{\gamma(x)}{1-\gamma(x)} \approx\ & 0.20 \pm 0.10 \quad \text{(SLAC)} \\ & 0.03 \pm 0.10 \quad \text{(EMC) (Drees and Montgomery, 1983),} \end{aligned} \tag{1.38}$$

where the larger experimental value is obtained for $Q^2 \sim 10\ \mathrm{GeV}^2$ typical of SLAC. This $R = 0.2$ corresponds to a fraction of scalar partons $\gamma \lesssim 0.25$. So most charged partons have half integer spin. The physical picture to explain a vanishing longitudinal cross-section for quarks is easy to understand. In a head on collision at high energies the helicity of the quark is conserved, i.e.



before

after

This can only occur if the helicity of the photon is $\lambda = \pm 1$, i.e. for transverse photons. On the other hand consider a spin-0-parton, then only a photon with helicity $\lambda = 0$ can be absorbed. Allowing for an initial transverse momentum and finite mass μ of the spin-1/2-parton one obtains $R = 4(k_T^2 + \mu^2)/Q^2$. In QCD (Reya, 1981) R becomes of order α_s, the QCD-coupling constant, due to the recoil of the quark having emitted a gluon. Unfortunately, up to now a reliable experimental determination of R does not exist (c.f. for reasons Drees and Montgomery (1983)).

Commonly one differentiates between different flavors in the quark distribution functions and their charges

$N_i(x)$:	$u(x)$	$c(x)$	$t(x)$	$d(x)$	$s(x)$	$b(x)$	$\bar{u}(x)$...
e_i :	2/3	2/3	2/3?	-1/3	-1/3	-1/3	-2/3...

The antiquark distributions are denoted by $\bar{u}(x)$, $\bar{d}(x)$, etc. For the proton the following normalizations have to be satisfied

$$\int[u(x) - \bar{u}(x)]dx = 2$$
$$\int[d(x) - \bar{d}(x)]dx = 1 \qquad (1.39)$$
$$\int[s(x) - \bar{s}(x)]dx = 0 \qquad \text{etc.}$$

$F_2(x)$ is given by correctly weighting the individual quark contributions, i.e.

$$F_2^p(x) = x\,\{\,\tfrac{4}{9}\,[u(x) + \bar{u}(x)] + \tfrac{1}{9}\,[d(x) + \bar{d}(x)] + \tfrac{1}{9}\,[s(x) + \bar{s}(x)] + \ldots\}. \qquad (1.40)$$

In neutrino reactions a charged $W^{+/-}$ is exchanged between the leptons and hadrons. Again there are three helicity states of the weak current and one can define cross-sections $\sigma(\lambda = +1)$, $\sigma(\lambda = -1)$ and $\sigma(\lambda = 0)$. Note, however, the current correlation function contains in

$$J_\mu^{weak}(x)\, J_\nu^{weak}(0) = \{V_\mu(x) - A_\mu(x)\}\{V_\nu(0) - A_\nu(0)\}, \qquad (1.41)$$

terms which are parity violating (-AV-VA); and two contributions (VV+AA) which are parity conserving. For a derivation of the cross-section of inclusive ν and $\bar{\nu}$-scattering we refer to the literature (Close, 1979), it is (M = target mass)

$$\frac{\partial^2\sigma(\overset{\nu}{\bar{\nu}})}{\partial\Omega\partial E'} = \frac{G^2}{2\pi^2}\, E'^2 \left(\frac{m_W^2}{m_W^2 + Q^2}\right)^2 \cdot$$

$$\cdot\,\{2W_1 \sin^2\theta/2 + W_2 \cos^2\theta/2 \pm W_3\, \frac{(E+E')}{M} \sin^2\theta/2\,. \qquad (1.42)$$

The photon exchange in eq. (1.42) $\frac{\alpha}{Q^2}$ has been replaced by $G = 1.2\cdot10^{-5}\ \mathrm{GeV}^{-2}$, the weak Fermi coupling constant[+], the extra factor $m_W^2/(m_W^2 + Q^2)$ corrects for the finite W-mass and can be ignored for all existing experiments. W_1 and W_2 parametrize the parity conserving interaction. W_3 arises from the interference terms -AV-VA. For the proton with mass $M = m$, $W_{1,2,3}$ are related to the parton distributions according to eqs. (1.43). The neutrinos measure down (d) and anti-up ($\bar{u}$) distributions and vice versa for the anti-neutrinos.

$$2m\,W_1^\nu = 2[d(x) + \bar{u}(x)] = F_1^\nu(x),$$
$$\nu W_2^\nu = xF_1^\nu(x) = F_2^\nu(x), \qquad (1.43)$$
$$\nu W_3^\nu = 2[d(x) - \bar{u}(x)] = F_3^\nu(x).$$

The sign difference in νW_3^ν in eqs. (1.42) and (1.43) concerning particle anti-particle conjugation is related to the change in the relative sign of the V and A terms in the weak current, when one goes from particles to anti-particles

[+]We set $\cos\theta_{Cabibbo} = 1$.

$$L^{\nu} = \bar{\mu}\gamma^{\nu}(1-\gamma_5)\nu_{\mu}$$
$$L^{\nu +} = \bar{\nu}_{\mu}\gamma^{0}[\gamma^{\nu}(1-\gamma_5)]^{+}\gamma^{0}\mu = \bar{\nu}_{\mu}\gamma^{\nu}(1+\gamma_5)\mu . \tag{1.44}$$

For anti-neutrinos we have from eq. (1.43)

$$\begin{aligned} 2m\, W_1^{\bar{\nu}} &= 2[u(x) + \bar{d}(x)] = F_1^{\bar{\nu}}(x), \\ \nu W_1^{\bar{\nu}} &= xF_1^{\bar{\nu}}(x) = F_2^{\bar{\nu}}(x), \\ \nu W_3^{\bar{\nu}} &= 2[u(x) - \bar{d}(x)] = F_3^{\bar{\nu}}(x). \end{aligned} \tag{1.45}$$

The fraction of momentum carried by anti-quarks in the proton can be extracted from the ratio

$$B = \frac{\int x(F_3^{\bar{\nu}} + F_3^{\nu})dx}{\int (F_2^{\bar{\nu}} + F_2^{\nu})dx} = \frac{2\int x[d(x) + u(x) - \bar{u}(x) - \bar{d}(x)]dx}{\int x[u(x) + d(x) + \bar{u}(x) + \bar{d}(x)]dx} . \tag{1.46}$$

One obtains for $< x >_{\bar{q}} \approx 0.05$. The anti-quarks carry only 5% of the momentum of the proton. Therefore one attributes them to the sea of soft partons. How much of the momentum do the quarks carry? An isoscalar target with baryon number 1 can be represented as the average of a free proton and neutron structure function. This average is denoted with N, then

$$\int dx\, x(u + \bar{u} + d + \bar{d} + s + \bar{s}) = \int dx(9\, F_2^{eN} - \frac{3}{2} F_2^{\nu N}) \approx 0.5. \tag{1.47}$$

So only fifty percent of the momentum is carried by quarks and anti-quarks. The residual fifty percent have to be attributed to gluons, the uncharged colored vector bosons mediating the strong interaction. Looking at the nucleon with this number in mind one would expect that also at rest about 50% of the energy+ is quark energy, the rest is contained in the bag or the confinement. This division goes well with string-like confinement, where the kinetic energy of the quarks N/R and the potential energy σR add up to the total energy

$$E_{tot} = N/R + \sigma R. \tag{1.48}$$

Minimizing the energy $\partial E/\partial R = 0$ makes both contributions equal. Substituting typical values for $N \approx 3x_0 \approx 6$ and $\sigma = 1$ GeV/fm one gets for the optimal size $R \approx 1$ fm. Unfortunately, the MIT-Bag-Model (Chodos and others, 1974) and several similar models give a bag energy $4\pi/3(B\cdot R^3)$ of the volume type. Thereby the ratio of quark to bag energy becomes unbalanced: quarks/bag = 3/1. Topological solitons (Kahana, Ripka and Soni, 1983) have the nice feature of giving an energy formula of the kind of eq. (1.48), unfortunately with constituents which have $(q\bar{q})$ substructure like pions. Probably such strong-binding solutions can be ruled out from deep inelastic scattering.

Güttner and others (1984) have analyzed the longitudinal π^+ electroproduction data $e + p \rightarrow e + n + \pi^+$ in terms of a pion distribution function of the proton. This interpretation is based on the assumption that the pion can be considered as a parton in the nucleon for low momentum transfer Q^2. We found a total percentage of 3% ± 0.5% $(\pi^+ n)$ in the proton. Adding up the other pionic components $(\pi^0 p, \pi\Delta, \pi N^*)$ we obtain as an upper limit less than 8% pionic content. The amount of momentum carried by the pions would be ≈ 1.6%. Consequently the nucleon has to be described

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+Using light-cone-coordinates xP^+ becomes x·mass at rest and we obtain the naive extrapolation above.

in terms of gluons and quarks. Nevertheless it is interesting to look at $(q\bar{q})$ correlations at low Q^2. This domain can be analyzed very well with the SURA machine. It is well known that the naive parton model with Bjorken scaling has to be corrected for low Q^2 and very high Q^2. At low $Q^2 \leq 5$ GeV^2 correlations between quarks are important in the same way as the two-body nucleon density $\rho_{ij}(k_i,k_j,q)$ is relevant for low q^2 nuclear structure functions. Final state interactions (Fig. 1.4) may also play a role at low Q^2. The Bag Model (Chodos and others, 1974) should be useful to calculate power corrections $1/Q^2$ to deep inelastic scattering (Figs. 1.4a, 1.4b). If the momentum transfer is small, the time $\Delta\tau \sim 1/Q$ is sufficiently long so that the struck quark can interact with a quark or anti-quark in its neighbourhood. These power corrections $O(1/Q^2)$ ("higher twist" corrections) have caused some problems in determining the parameter Λ_{QCD} related to scaling violations at large Q^2. In the limit of $Q^2 \gg \Lambda^2_{QCD}$ the strong interactions between partons become weak and can be treated perturbatively. The deviations from naive scaling due to QCD are one of the cornerstones to establish Quantum Chromodynamics as the theory of strong interactions (Figs. 1.4c, 1.4d). Deep inelastic lepton-nucleus scattering at large Q^2 has to take these effects into account, too. Therefore let us discuss them now; for more detailed reviews see Buras (1981); Nachtmann (1980); Pennington (1983); Reya (1981).

A photon with invariant mass Q^2 resolves distances $\Delta r_\perp \sim 1/Q$ and (light-cone) times $\Delta\tau \sim 1/Q$. The interpretation of the deep inelastic scattering in terms of distribution functions $N_i(x,Q^2)$ can be maintained, when one distinguishes the "size" of the constituents, which the photon can resolve. At low Q^2 a quark and

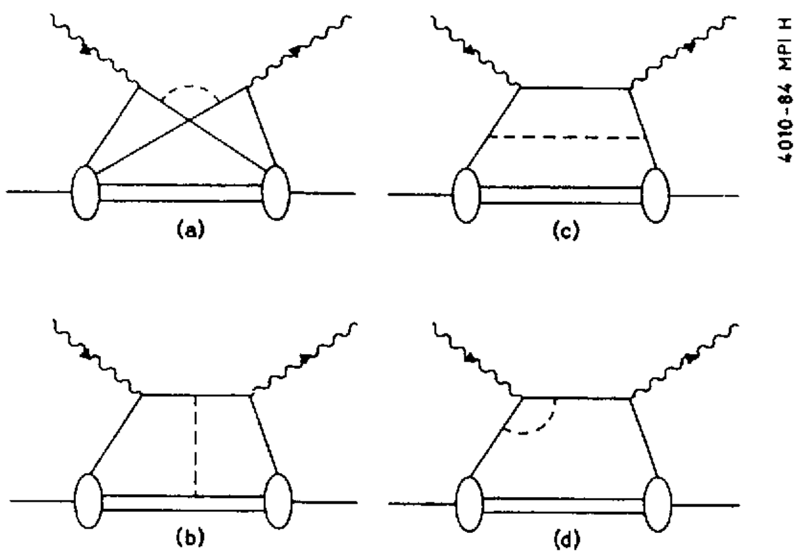

Fig. 1.4. Correlations (1.4a) and final state interactions (1.4b) important at low Q^2 ("twist-4 contributions"). Scaling corrections (1.4c, 1.4d) at large Q^2 ("twist-2 effects"). Dotted lines are gluons.

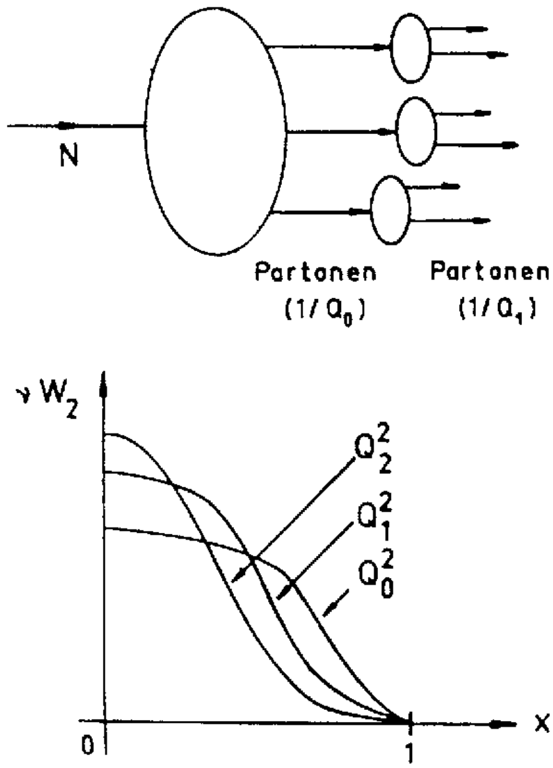

Fig. 1.5. Softening of the x-distribution for "smaller" ($1/Q_2 < 1/Q_1 < 1/Q_0$) constituents.

gluon nearby at $\Delta r_\perp \leq 1/Q$ look like one quark with a "size" $1/Q$. We do not have to calculate their interaction, they form one effective quark. We treat deep inelastic scattering in impulse approximation for this effective constituent of size $1/Q$. Increasing the resolution of the photon, however, resolves the composite quark gluon system. It increases the number of partons, which will manifest itself qualitatively as a softening of the structure function because more "smaller" partons have a smaller average momentum fraction each (Fig. 1.5). The amount of sea quarks increases, whereas the valence quarks lose momentum due to radiating soft gluons.

We will take the dependence of the distribution functions on the effective quark size $1/Q^2$ into account by defining $N_i(x,Q^2)$ as the probability to find a quark of size $1/Q^2$ with momentum fraction x. We can keep the relation between the structure function νW_2 and the quark distribution function, i.e.

$$\nu W_2(\nu,Q^2) = \sum_i e_i{}^2 x \cdot N_i(x,Q^2). \tag{1.49}$$

How do the quark and gluon distribution functions change with Q^2? Assume the parton distribution $N_i(x,Q_0^2)$ is known and we want to calculate $N_i(x,Q_1^2)$ with

$$Q_1^2 > Q_0^2. \tag{1.50}$$

Due to the increase on the cut-off of transverse momenta from Q_0^2 to Q_1^2 new modes of the QCD-fields, i.e. new quark and gluon modes, can become occupied.

The gluon distribution $\Delta N_G(x)$ resulting from the bremsstrahlung of a quark having momentum fraction $y = 1$ is obtained from first-order perturbation theory (Fig. 1.6)

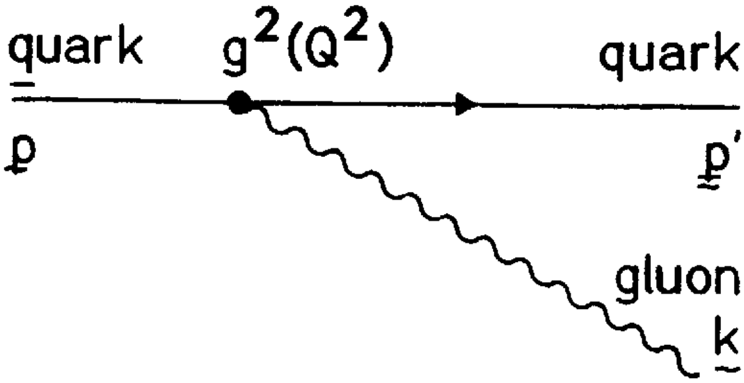

Fig. 1.6. Quark radiating a gluon. The gluon mode is only allowed in the transverse space $Q_0^2 \leq k_\perp^2 \leq Q_1^2$, which is resolved by increasing the photon momentum from Q_0^2 to Q_1^2.

$$\Delta N_G(x) = \int_{Q_0^2 \leq k_\perp^2 \leq Q_1^2} \frac{d^3k}{(2\pi)^3} \frac{1}{k_0} \frac{d^3p'}{(2\pi)^3} \frac{1}{p_0'} \delta^3(\vec{p} - \vec{p}' - \vec{k})\delta(x - \frac{\vec{k}\cdot\vec{p}}{p^2}) \cdot$$

$$\cdot \left| \frac{< G^a(k)q(p')| H_{QCD}| q(p) >}{|\vec{k}| + |\vec{p}'| - |\vec{p}|} \right|^2 = \tag{1.51}$$

$$= \frac{g^2(\bar{Q}^2)}{8\pi^2} \frac{4}{3} \frac{1 + (1 - x)^2}{x} \ln \frac{Q_1^2}{Q_0^2} = \frac{g^2(\bar{Q}^2)}{8\pi^2} P_{Gq}(x)\Delta \ln Q^2.$$

One recognizes the familiar logarithmic dependence of bremsstrahlung weighted with the running coupling constant of QCD $g^2(Q^2)$, which at large Q^2 has the dependence (Wilczek , 1982; Reya, 1981)

$$\alpha_s(Q^2) = \frac{g^2(Q^2)}{4\pi} = \frac{4\pi}{(11 - 2f/3)\ln Q^2/\Lambda^2} , \tag{1.52}$$

where the scale parameter of QCD is $\Lambda \approx 100 - 300$ MeV and f is the number of flavors important in the problem. For a general initial distribution of partons $N_i(y,Q^2)$ the Altarelli-Parisi equation summarizes the Q^2-evolution of the distribution functions. Generalizing eq. (1.51) one obtains (Altarelli and Parisi, 1977)

$$\frac{\partial N_i(x,Q^2)}{\partial \ln Q^2} = \frac{g^2(Q^2)}{8\pi^2} \int_0^1 dy \int_0^1 dz\; P_{i\leftarrow j}(z)\; N_j(y,Q^2)\; \delta(x - zy), \tag{1.53}$$

where i and j mean quarks or gluons and $P_{ij}(z)$ is the probability of a constituent j to emit a constituent i with momentum fraction z. $P_{i=\text{gluon},j=\text{quark}}$ has been given in eq. (1.51); for others consult the literature (Reya, 1981; Pennington, 1983). Since the folding in eq. (1.53) is multiplicative, the moments of the structure functions obey first-order differential equations. The moments of the non-singlet distribution function $N^{NS}(x,Q^2)$ behave particularly simple. How is the non-singlet function defined? It corresponds to the average structure function when neutron and proton are in an I = 1 state. Similarly the average I = 0 deuteron structure function defines the singlet distribution for flavor SU2. For the proton we have eq. (1.40) and for the neutron we use isospin invariance, which makes the distribution of u(n) and d(n)-quarks in the neutron equal to the distribution of d and u-quarks in the proton, i.e. u(n) = d(p) = d, d(n) = u(p) = u.

$$\begin{aligned} F_2^{ep}/x &= \frac{4}{9}\,(u+\bar{u}) + \frac{1}{9}\,(d+\bar{d}) + \ldots = \frac{5}{18}\,(u+\bar{u}+d+\bar{d}) + \frac{1}{6}\,(u+\bar{u}-d-\bar{d}) \\ F_2^{en}/x &= \frac{1}{9}\,(u+\bar{u}) + \frac{4}{9}\,(d+\bar{d}) + \ldots = \frac{5}{18}\,(u+\bar{u}+d+\bar{d}) - \frac{1}{6}\,(u+\bar{u}-d-\bar{d}). \end{aligned} \tag{1.54}$$

Consequently the non-singlet structure function corresponds to the difference

$$N^{NS}(x,Q^2) = (F_2^{ep} - F_2^{en})/x. \tag{1.55}$$

The Altarelli-Parisi equation for $N^{NS}(x,Q^2)$ does not differentiate between quarks and anti-quarks. Since it is an equation for probability densities, it contains transition probabilities P_{ij}, which are the same for quarks and anti-quarks radiating gluons. Therefore the coupling between gluons and quarks drops out for the non-singlet structure function.+ Using eqs. (1.52) and (1.53) one obtains for the moments of N^{NS}

$$M_n^{NS}(Q^2) = \int_0^1 x^{n-1}\, N^{NS}(x,Q^2)\,dx \tag{1.56}$$

the following equation

$$\frac{d\, M_n^{NS}(Q^2)}{d \ln Q^2} = - \frac{1}{\ln Q^2/\Lambda^2}\, d_{qq}^n\, M_n^{NS}(Q^2) \tag{1.57}$$

with the anomalous dimension d_{qq}^n

$$d_{qq}^n = \frac{-6}{(33 - 2f)} \int dz\, z^{n-1}\, P_{qq}(z). \tag{1.58}$$

The solution of eq. (1.57) $(M_n^{NS})^{-1/d_{qq}^n}$ must depend linearly on ln Q^2. The anomalous

+The same holds for the average neutrino structure function $F_3^{\nu N} = \frac{1}{2}\,(F_3^{\nu p} + F_3^{\nu n})$.

dimensions d^n_{qq} have a direct interpretation in the field theoretic treatment of deep inelastic scattering where one makes a taylor expansion of the bilocal operator product $j_\mu(x) j_\nu(0)$ (Reya, 1981).

Let us summarize this chapter on the parton model of the nucleon. In the infinite momentum frame the nucleon can be approximated as a jet of partons which interact with the virtual photon as point charges. This leads to Bjorken scaling, i.e. a dependence of the structure function $\nu W_2(\nu, Q^2)$ and $2m\, W_1(\nu, Q^2)$ on $x = Q^2/2m$ only. There are weak violations of Bjorken scaling due to QCD-radiation.

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II. New Topics in Deep Inelastic Lepton Nucleon Scattering

1. Where is the proton spin?

One of the spectacular results in deep inelastic lepton nucleon scattering during the last year has been the polarized spin structure function of the proton. By scattering polarized muons on a polarized hydrogene target, one can measure the assymmetry A of the difference of two cross sections $d\sigma\uparrow\downarrow - d\sigma\uparrow\uparrow$ over their sum Using known kinematical

$$A = \frac{d\sigma^{\uparrow\downarrow} - d\sigma^{\uparrow\uparrow}}{d\sigma^{\uparrow\downarrow} + d\sigma^{\uparrow\uparrow}} = D\,(A_1 + \eta A_2) \approx D A_1$$

factors one can deduce the difference of quark distributions of the species i with quark helicity parallel and antiparallel to the proton spin. The quantization axis is the spin of the proton.

$$A_1 = \frac{\sum e_i^2 \left(f_i^{\uparrow}(x) - f_i^{\downarrow}(x) \right)}{\sum e_i^2 \left(f_i^{\uparrow}(x) + f_i^{\downarrow}(x) \right)}$$

The x dependence of this difference relative to the sum is shown in figure 1. There are currently two proposals to measure the same quantity on the neutron with a polarized deuteron target or with a deuteron jet target.Let us first discuss the integral over this difference which can be related to the total proton spin.

$$\int_0^1 dx \, \frac{A_1(x) \, F_2(x)}{2x(1+R)} = \int dx \, g_1(x) = 0.122 \pm 0.013$$

This quantity has attracted most of the attention. Why is it so exciting? Using the known measurements of the neutron β-decay and hyperon -decays one con calculate from

$$\Delta u - \Delta d = 1.23 \pm 0.1$$

$$\Delta u + \Delta d - 2\Delta s = 0.68 \pm 0.01$$

the measured integral the flavor singlet sum of quarks with spin parallel minus antiparallel to the proton spin.

$$\Delta q = \int_0^1 \left\{ \left[q\uparrow(x) - q\downarrow(x) \right] + \left[\bar{q}\uparrow(x) - \bar{q}\downarrow(x) \right] \right\} dx$$

$$\begin{aligned} 0.23 = 2 \int_0^1 g_1(x)\, dx = \ & \frac{1}{6} \left(\Delta u - \Delta d \right) \\ & + \frac{1}{18} \left(\Delta u + \Delta d - 2\Delta s \right) \\ & + \frac{2}{9} \left(\Delta u + \Delta d + \Delta s \right). \end{aligned}$$

$$\Delta u + \Delta d + \Delta s = 0.00 \pm 0.02$$

The result is zero, which is very surprising since we would expect in the constituent quark model that there is one more quark with spin parallel to the proton spin relative to the antiparallel ones. The experimental result of zero for the same quantity prompted the question: Where is the spin of the proton? As far as I can see there is not yet any convincing answer to this naive question. Various theoretical ideas are circulating with partial insight into this delemma. It is well known that there exists a problem with the singlet axial current. Its divergence is not equal to the pseudoscalar quark density times twice the quark mass as for free quarks but there exists an extra piece which comes from the gluons. It expresses the product of $E.B$ which is a pseudoscalar. The most clear derivation of this property is given by Fujikawa[1]. It is based on the fact that under chiral singlet rotations the fermion integration measure in gauge theories acquires a phase due to different behaviour of right and lefthanded fermions in an external vector field. The classical spin density is defined as the the spin density the divergence of which gives the classical pseudoscalar operator. It then represents the measured quark spin density which equals the quark spin density minus the gluon spin density. The measured result would imply that the gluon spin compensates the quark spin. Unfortunately none of the pictures of the proton we have gives any hint about this gluon spin. A more heuristic relation has been discussed by Veneziano[2] and by Hatsuda [3]. It relates the singlet pseudoscalar coupling G_A to the spin density. A natural consquence of the vanishing of the integral $\Delta u + \Delta d + \Delta s$ is then the vanishing of the η' -coupling to the nucleon. Note it is the heavy

η' which has the singlet SU(3) property. I personally find the x dependence of the function

$$g(\eta' NN) f_\pi \sqrt{2N_f} = 2m_N (\Delta u + \Delta d + \Delta s) = 2m_N G_A^{(1)} \approx 0$$

$g_1(x)$ as interesting as the integral. Our complete inability to calculate the properties of the proton manifests itself in this quantity in clear fashion. The constituent model gives no x-dependence to the function g_1 , which would be equal to 5/9. Where does the x-dependence come from? For the large x-region we know that the proton wavefunction does not have the usual amount of u to d-quarks. Nature seems to prefer the u-quark at large x. This can be seen from the ratio of $F_2(x)$ for the neutron to the proton. A simple estimate gives the empirical value if u-quarks dominate. One can speculate that the up quark coupled to the ud pair with I=0 and spin=0 combination in the proton wavefunction dominates at large x .This residual up quark would then carry the proton isospin and the proton spin giving the increase of g_1 to 0.7.

At low x the radiation of gluons will produce as many sea quark pairs with spin parallel as antiparallel, whereas the parent parton has a preferred spin which it guards. In this case the spin density of the constituent quark model is diluted and g_1 goes to zero.

2 – Shadowing and Antishadowing-New Aspects of the EMC-Effect

The EMC effect was discovered in 1983. Since then a huge experimental effort has been devoted to measure the nuclear structure function over a larger domain in the Bjorken variable x. Especially the low x-region is now vastly extended compared to the situation in

1983 [7], see fig.2. The current -current correlation function $W_{\mu\nu}$ is light cone dominated

$$W_{\mu\nu}(\nu, Q^2) = \frac{1}{4\pi} \int d^4y \, e^{-iqy} \langle p | [J_\mu(y), J_\nu(0)] | p \rangle$$

as one can see from the kinematics in the rest frame of the hadron. For fixed Bjorken $x = Q^2/(2M\nu)$ the energy transfer ν becomes $-Mx + |q| \rightarrow |q|$. Therefore the exponent of the Fouriertransform selects distances $y_0 - y_z = 0$ which correspond to correlations on the light cone. The quark times antiquark field operator in the middle of $W_{\mu\nu}$ can be replaced by the free propagator singularity. The symmetric Lorentz tensor in μ and ν gives the probability of annihilating a quark a position 0 and time 0 and recreating it at $+y_z$ and time y_0. From the arguments of the exponential one sees that the distances involved are of the order of $1/Mx$ in $y_0 + y_z$ whereas both $y_0$and y_z are on the light cone i.e. $y_0 = y_z$. This simple algebra makes obvious that in the shadowing region we test quark correlations over long distances in the nucleus. E.g. for $x = 0.05$ we find the distance $d = 4fm$. In virtual photon nucleus interactions a quark antiquark system propagates as a virtual state of the photon through the nucleus the transverse dimension of which can be estimated from the transverse momentum over the longitudinal momentum multiplied with the characteristic longitudinal distance $1/Mx$. Here we take a transverse momentum of the order of the

$$r_\perp = \left| \frac{\vec{p}_\perp}{p_z} \right| \cdot \frac{1}{Mx} = \frac{m_\pi}{\sqrt{\nu^2 + Q^2}} \, \frac{1}{Mx} = \frac{m_\pi}{Mx\,\nu} \approx \frac{2m_\pi}{Q^2}$$

pion mass. If the size of this state is of the size of a hadron, we see shadowing under the condition that the nuclear radius is much larger than the mean free path which is of the order of $2 - 3fm$ for a hadron. As shown above the radial size of the q and $antiq$ system decreases with increasing Q^2, thereforeshadowing will decrease with increasing Q^2. Both of these effects are nicely shown in a calculation based on the generalized vecot meson dominance model by Schildknecht et al [4]in figure 3. A very interesting physical picture has been proposed also recently by Mueller [5] and worked out by Qiu and Close[6]. In the infinite momentum frame of deep inelasic scattering the nucleus is Lorentz contracted. Also the valence quarks in the nucleons inside the nucleus become thin as pancakes. The sea quarks , however, have an extension which depends on a cut-off in x space which normally is taken as $x_{cut} = p_{tr}^2/\Lambda^2$ or $x_{cut} = p_{tr}^2/ Q^2$. It is invariant under Lorentz boosts. Therefore at a given γ-factor of the boost the extension of the sea quarks of each individual nucleon will be of the size of the whole contracted nucleus, The overcrowding of the sea quarks leads to merging of the quarks and anti- quarks among them and therefore diminishes the deep inelastic yield, we see shadowing. Up to now the model is not yet quantitive but it promises to discuss the EMC effect in the low x-region in the same way as in the intermediate x-region.

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Figure Captions

Figure 1. The spin structure function A_1 of the proton as measured by the EMC-Collaboration [8].

Figure 2. The ratio of structure functions for Ca over deuteron, C/d and Ca/C from ref [7]

Figure 3. Calculation of shadowing from ref [4] .Data from ref [7]

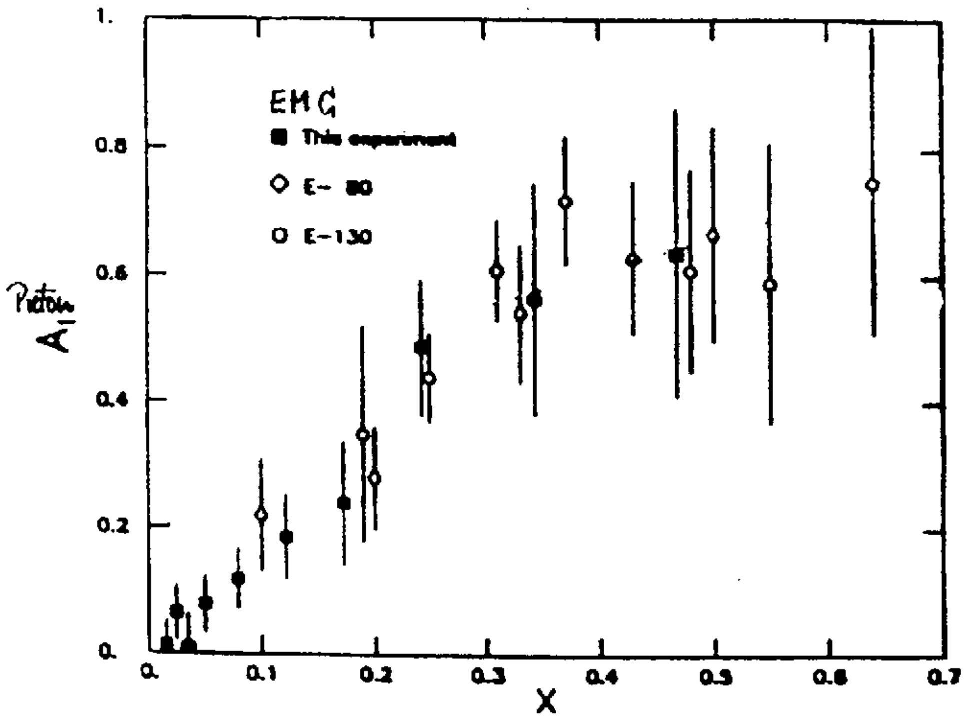
EMC
This experiment
E- 80
E-130
Proton
A_1
X
0.
0.2
0.4
0.6
0.8
1.
0.1
0.3
0.5
0.7

fig. 1

C. Scholz
Thesis
NMC-
Results

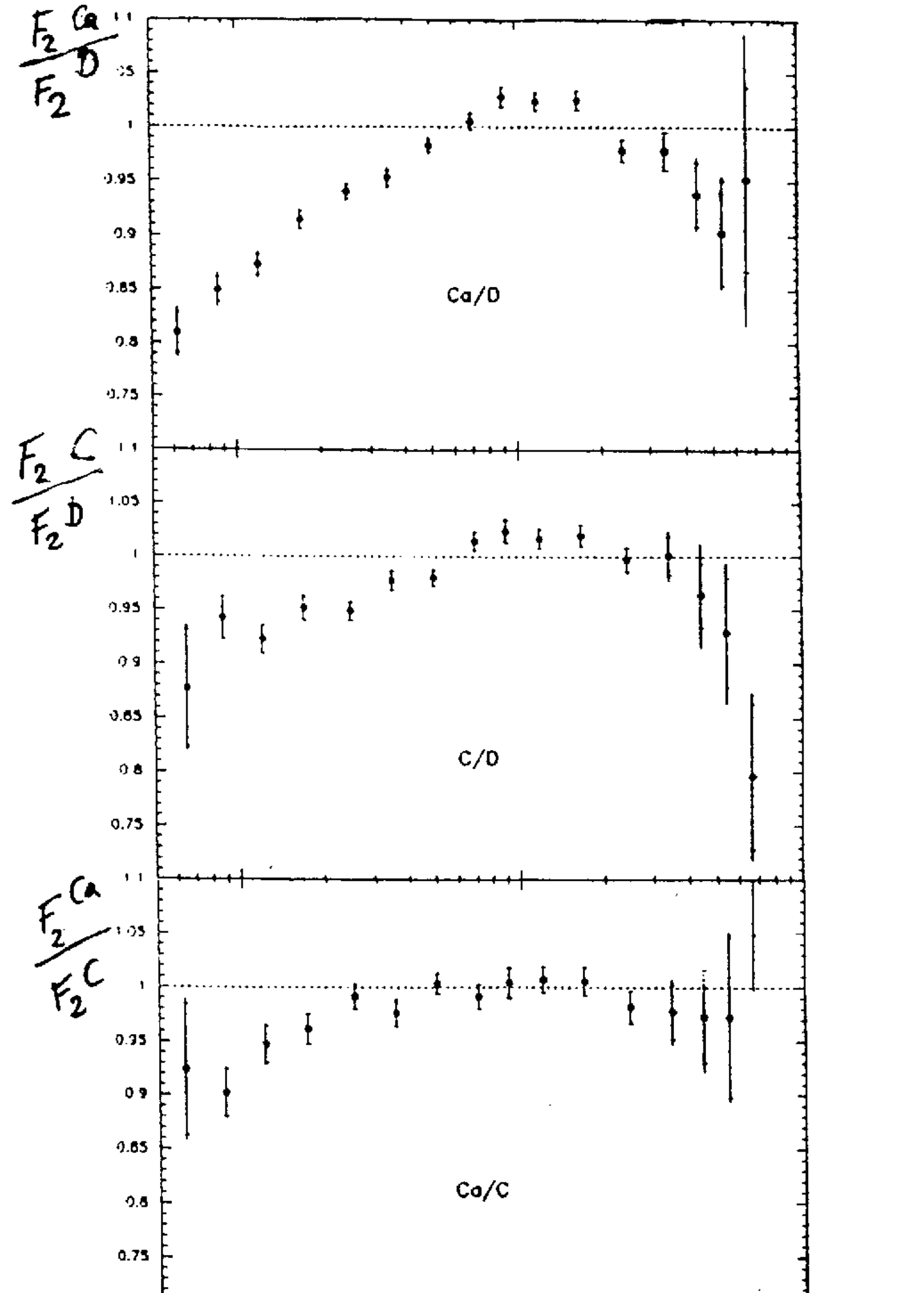

fig. 2