Nuclear Effects in Structure Functions

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Abstract

A review of the experimental and theoretical results in the field of nuclear effects in deep inelastic scattering is presented.

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Acknowledgements

Chapter 1

Introduction

Charged lepton Deep Inelastic Scattering (DIS) off nucleons has been for a long time a powerful tool to investigate the structure of matter.

In the late 1960's "scaling" of the nucleon structure functions was discovered by the pioneering SLAC experiments. Scaling was interpreted as evidence that nucleons contain pointlike constituents – later identified with quarks. About ten years later "scaling violations" were detected, again by deep inelastic scattering experiments; this showed that quarks interact as predicted by the theory of strong interactions, Quantum Chromo-Dynamics (QCD).

Nuclear effects in deep inelastic scattering were thought to be largely negligible, essentially because the energies involved in deep inelastic scattering are much larger than those typical of nuclear processes. The nucleus was viewed as a collection of quasi-free nucleons and quarks were believed to be insensitive to the nuclear environment. This picture was however proven wrong by a comparison of the iron and deuterium structure functions, measured by the European Muon Collaboration at CERN in 1982. The phenomenon has since been actively studied, both experimentally and theoretically; it is our purpose to review these efforts.

Much of the interest in this subject comes from the fact that its understanding may help to bridge the gap between nuclear and particle physics: QCD has been successful in describing a number of short distance phenomena, at scales $\ll 1$ fm; unfortunately, at the moment, it appears very difficult to derive from this theory quantitative predictions for processes involving larger distances, such as those typical of nuclear physics. The study of the nuclear effects in deep inelastic scattering may shed some light on the large distance behavior of quarks and gluons.

1.1 Generalities of Charged Lepton Deep Inelastic Scattering

This section is meant to give a brief introduction to deep inelastic scattering, mainly intended to define some basic quantities and concepts; a more detailed presentation of the subject can be found elsewhere (see e.g. [1]).

1.1.1 Variables and Cross Section

We are interested in the process $l + N \rightarrow l + X$, where a charged lepton scatters inelastically off a nucleon; the final state contains the scattered lepton and the debris of the nucleon. We consider the case in which the target nucleon is at rest and we assume that only the incoming and outgoing lepton momenta are measured. The Feynman graph for the process is shown in fig. 1.1, under the assumption that only one virtual photon is exchanged.

If E and E' are the energies of the incident and scattered lepton, k, k' their four momenta, and ϑ is the lepton scattering angle, the following kinematic variables can be introduced (we neglect the lepton mass):

- $q = k - k'$, the four momentum of the virtual photon,
- $Q^2 = -q^2 = -(k - k')^2 \simeq 4EE' \sin^2(\vartheta/2)$, with q^2 the square of the four momentum transferred from the lepton to the target nucleon,
- $\nu = P \cdot q/M = E - E'$, the virtual photon energy,
- $W^2 = (P + q)^2 = M^2 - Q^2 + 2M\nu$, the total invariant mass squared of the target nucleon-virtual photon system,
- $x = Q^2/(2P \cdot q) = Q^2/(2M\nu)$, the Bjorken scaling variable,
- $y = P \cdot q/(P \cdot k) = \nu/E$,

where P is the four momentum of the target nucleon and the rightmost expressions for Q^2 , ν , x and y are valid in the laboratory frame, where the target nucleon is at rest. The mass M is conventionally taken to be that of the proton.

The expression for the cross section of this process reads

$$\frac{d^2\sigma}{d\Omega dE'} = \frac{4\alpha^2 E'^2}{Q^4} \left(2W_1 \sin^2 \frac{\vartheta}{2} + W_2 \cos^2 \frac{\vartheta}{2} \right), \quad (1.1)$$

with α the fine structure constant. The structure functions $W_1(\nu, Q^2)$ and $W_2(\nu, Q^2)$ contain information about the inner structure of the target nucleon; they can be related to the photo-absorption cross sections σ_T and σ_L for transverse (i.e. with helicity ± 1) and longitudinal photons (helicity 0):

$$\begin{aligned} W_1 &= \frac{K}{4\pi^2\alpha} \sigma_T \\ W_2 &= \frac{K}{4\pi^2\alpha} (\sigma_T + \sigma_L) \frac{Q^2}{Q^2 + \nu^2}, \end{aligned}$$

where K is the incident photon flux. For virtual photons there is some arbitrariness in the definition of the photon flux: Gilman's convention is to take $K = \sqrt{\nu^2 + Q^2}$; alternatively one can have (Hand's convention) $K = \nu - Q^2/(2M)$. Both reduce to $K = \nu$ for real photons ($Q^2=0$).

The longitudinal to transverse virtual photon cross section ratio

$$R = \frac{\sigma_L}{\sigma_T} = \frac{W_2(\nu, Q^2)}{W_1(\nu, Q^2)} \left(1 + \frac{\nu^2}{Q^2} \right) - 1 \quad (1.2)$$

is an important quantity, being related (see below) to the spin of the nucleon's constituents.

In terms of σ_L and σ_T the cross section 1.1 becomes

$$\frac{d^2\sigma}{d\Omega dE'} = \Gamma(\sigma_T + \epsilon\sigma_L),$$

where

$$\epsilon = \left(1 + 2 \frac{Q^2 + \nu^2}{Q^2} \tan^2 \frac{\vartheta}{2} \right)^{-1},$$

and

$$\Gamma = \frac{K \alpha}{2\pi^2 Q^2} \frac{E'}{E} \frac{1}{1 - \epsilon}.$$

It is customary to define

$$\begin{aligned} F_1 &= MW_1 \\ F_2 &= \nu W_2 \end{aligned}$$

and rewrite 1.1 in terms of x and Q^2 :

$$\frac{d^2\sigma}{dx dQ^2} = \frac{4\pi\alpha^2}{Q^4} \left[\left(1 - y - \frac{Mxy}{2E} \right) \frac{F_2(x, Q^2)}{x} + y^2 F_1(x, Q^2) \right]. \quad (1.3)$$

Often R is used instead of F_1 , in which case we have:

$$\frac{d^2\sigma}{dx dQ^2} = \frac{4\pi\alpha^2}{Q^4} \frac{F_2(x, Q^2)}{x} \left[1 - y - \frac{xyM}{2E} + \frac{y^2}{2} \frac{1 + 4M^2x^2/Q^2}{1 + R(x, Q^2)} \right]. \quad (1.4)$$

1.1.2 Scaling and the Quark-Parton Model

The first measurements of the structure functions F_1 and F_2 were made with electron beams at SLAC and DESY at the end of the 60's; they were followed by extensive and detailed investigations both with electron and muon beams. Figure 1.2 shows a recent compilation of results on the deuterium structure function $F_2^D(x, Q^2)$ [2].

It is clear from fig. 1.2 that F_2 is to first approximation independent of Q^2 . In general both structure functions F_1 and F_2 appear to depend mainly on x ; moreover they tend to a finite value even when both Q^2 and ν are large, as long as the ratio Q^2/ν is kept constant, the so-called "Bjorken limit". This scaling behavior was unexpected and strongly contrasts with the properties of elastic form factors, which instead rapidly vanish with increasing Q^2 .

Scaling can be interpreted as a sign of the fact that the virtual photon is absorbed, inside the nucleon, by pointlike, electrically charged objects, first named partons by Feynman and later identified with quarks.

Indeed one can show that the observed cross section can arise from the incoherent sum of the cross sections for elastic scattering of the lepton off pointlike, spin-1/2, charged objects. In a frame where the momentum of the nucleon is very large (the "infinite momentum frame"), the parton on which the interaction takes place carries a fraction $x = Q^2/(2M\nu)$ of the nucleon's momentum. Of all possible infinite momentum frames the Breit frame is often used, in which the virtual photon has zero energy and the parton exactly reverses its longitudinal momentum

in the collision. The main result of this approach is that F_1 and F_2 can be expressed in terms of the distribution function of the partons inside the nucleon:

$$F_1 = MW_1(\nu, Q^2) = \frac{1}{2} \sum_i e_i^2 q_i(x), \quad (1.5)$$

where the sum is over all parton species, e_i is the charge of the i -th parton species and q is its distribution function, which – in this approach – depends on x only. Similarly:

$$F_2 = \nu W_2(\nu, Q^2) = x \sum_i e_i^2 q_i(x). \quad (1.6)$$

The structure function F_2 is thus the sum over the various parton species of the partons' momentum distributions $xq_i(x)$ weighted by their charges squared.

From expressions 1.5 and 1.6 the Callan-Gross relation can be derived:

$$F_2(x) = 2xF_1(x). \quad (1.7)$$

Finally we remark that formula 1.2 now reads:

$$R = \frac{4M^2 x^2}{Q^2}, \quad (1.8)$$

which tends to zero in the Bjorken limit.

The function R depends strongly on the spin of the partons. If one assumes spin-0 partons, instead of spin-1/2, one finds $R = \infty$; R is thus a sensitive indicator of the spin of the partons.

The identification of partons with the spin-1/2, fractional charge quarks first introduced by Gell-Mann and Zweig is now universally accepted. One normally refers to the Quark-Parton Model (QPM), in which nucleons are described as composed of *valence* quarks, which define the charge and spin of the nucleon. In the proton there are two quarks of “flavor” u and one of flavor d . In the neutron there are two d quarks and one u quark. The electric charge of u quarks is $(2/3)e$, that of d quarks is $-(1/3)e$. The QPM, in its generally accepted formulation, goes beyond this and assumes that quarks interact via the exchange of spin-1 massless particles, the gluons (see next section), which can fluctuate into quark-anti-quark pairs, generally referred to as *sea* quarks.

1.1.3 Scaling Violations and QCD

Close examination of fig. 1.2 shows that F_2 does have a weak dependence on Q^2 . Structure functions rise with increasing Q^2 for $x < 0.2$ and then start decreasing.

As we mentioned above, partons interact among themselves and their behavior is described by Quantum Chromo-Dynamics (QCD), a non-abelian gauge theory. According to QCD each quark flavor may carry one of three possible “color” charges; quarks interact by exchanging spin-1, massless, electrically neutral particles called gluons, which are themselves colored.

In analogy with Quantum Electro-Dynamics (QED), quarks may radiate gluons and gluons may convert into quark anti-quark pairs. Unlike what happens in QED however, gluons, being colored, may directly interact with each other.

The QCD coupling constant has the following form:

$$\alpha_s(Q^2) = \frac{12\pi}{(33 - 2N_f) \ln(Q^2/\Lambda^2)}, \quad (1.9)$$

where N_f is the number of quark flavors and Λ is a mass scale parameter, the only free parameter in the theory to be determined from experiment. It is evident that the value of α_s tends to zero when Q^2 tends to infinity. This remarkable property is known as asymptotic freedom and, for large Q^2 , justifies a posteriori the quark-parton model treatment, where quarks are taken as essentially free. Asymptotic freedom is also the basis for the application of perturbative methods to large Q^2 phenomena (perturbative QCD). Non-perturbative approaches, although actively investigated, imply formidable computational problems and are – to date – of scarce predictive power.

Perturbative QCD is unable to predict the shape of the structure functions at fixed Q^2 . It can however make predictions about the rate of change (the evolution) of the structure functions with Q^2 , as long as Q^2 is large enough, so that α_s is sufficiently small and perturbative calculations are meaningful. The evolution of structure functions may be investigated by studying the evolution of their moments $M_n^{(j)}$

$$M_n^{(1)} = \int_0^1 dx x^{n-1} F_1(x, Q^2), \quad (1.10)$$

$$M_n^{(2)} = \int_0^1 dx x^{n-2} F_2(x, Q^2), \quad (1.11)$$

or, more easily, by using the Altarelli-Parisi evolution equations:

$$\frac{dq_i(x, t)}{dt} = \frac{\alpha_s(Q^2)}{2\pi} \int_x^1 \frac{dx'}{x'} \left[\sum_i q_i(x', t) P_{qq}(x/x') + G(x', t) P_{qG}(x/x') \right] \quad (1.12)$$

$$\frac{dG_i(x, t)}{dt} = \frac{\alpha_s(Q^2)}{2\pi} \int_x^1 \frac{dx'}{x'} \left[\sum_i q_i(x', t) P_{Gq}(x/x') + G(x', t) P_{GG}(x/x') \right], \quad (1.13)$$

where

- the sums are over all quarks and anti-quarks;
- $t = \ln(Q^2/\mu^2)$ (μ^2 is a scale unspecified in perturbative QCD, see below);
- x' is the momentum fraction carried by the parton;
- $q_i(x', Q^2)$ is the distribution function of quarks of flavor i ;
- $G(x', Q^2)$ is the gluon distribution function;
- $P_{qq}(x/x')$, $P_{qG}(x/x')$, $P_{Gq}(x/x')$ and $P_{GG}(x/x')$ are the quark-quark, quark-gluon, gluon-quark and gluon-gluon splitting functions. $P_{ab}(x/x')$ measures the probability for finding a parton (we mean quarks and gluons, now) of type a with momentum fraction x “inside” (i.e. radiated from) a parton of type b and momentum fraction x' .

In all QCD formulae seen above (notably in 1.9 and 1.12-1.13) Q^2 always appears scaled either by Λ^2 or μ^2 . The magnitude of Λ controls the rate at which the QCD coupling constant runs as a function of Q^2 ; Λ is found to be ≈ 200 MeV. The scale μ can be regarded as the low-momentum cutoff for gluon radiation; in other words, for $Q^2 = \mu^2$ the nucleon appears simply as made up of three valence quarks. It can be argued that $\mu^2 \sim 1/\lambda^2$, where λ is the quark confinement size.

1.2 Nuclear Effects in DIS

Nuclear effects in DIS are studied by comparing the structure functions F_2 of bound nucleons and of deuterium, the latter being taken as the best approximation to the structure function of a free, isoscalar nucleon. Direct comparison with the structure function of the proton, strictly speaking the only free nucleon, is not very significant in view of the fact that nuclei contain protons and neutrons in approximately equal amounts and that the proton and the neutron structure functions are different, especially at large x .

Note that here and throughout this paper we always refer to F_2 *per nucleon* and not per nucleus; likewise cross sections measured on nuclei will always be intended per nucleon, that is after dividing the cross section per nucleus by the number of nucleons in the nucleus, A . Furthermore, unless otherwise stated, all results are corrected for the (generally small) non-isoscalarity of the targets: for instance the ^{56}Fe to D structure function ratio refers to an ideal Fe nucleus containing an equal amount of protons and neutrons.

The structure function measured on a bound nucleon, F_2^A , is found to differ from that measured on deuterium. Naively, keeping in mind expression 1.6 for F_2 in the QPM, this implies that the quark distribution functions are different for a bound and a free nucleon: somehow a quark “feels” if the nucleon to which it belongs is free or is embedded in a nucleus.

The fact that $F_2^A \neq F_2^D$ is referred to with different terms, depending on the region of x . Figure 1.3 shows the ratio F_2^A/F_2^D as a function of x in a schematic way; with reference to this figure we now introduce some of the terminology.

- For $x < 0.05-0.1$ $F_2^A < F_2^D$ and the ratio F_2^A/F_2^D is smaller than unity. This depletion is known as “shadowing”.
- For $x \approx 0.1-0.2$ the ratio is larger than unity by a few per cent; this enhancement is sometimes called “anti-shadowing”.
- At $x \approx 0.2-0.3$ the ratio starts decreasing; it reaches a minimum around $x = 0.6$ and then begins to rise again.

The behavior of the ratio in the region $0.1-0.2 < x < 0.8$ is often referred to as the “EMC effect”; this nomenclature is however not accepted universally and the term EMC effect is sometimes used to label all nuclear effects in structure functions.

- For $x > 0.8$ the ratio grows above unity; this is the “Fermi motion” region.

An important remark is in order. The structure function ratio F_2^A/F_2^D is identically equal to the cross section ratio σ^A/σ^D if the quantity $R = \sigma_L/\sigma_T$ does not depend on A . This can be seen by rewriting relation 1.4 for a nuclear target A and for D, and then taking the ratio. As we shall discuss, there is experimental evidence indicating that R is indeed A -independent, at least in a certain kinematic region.

Chapter 2

Experimental Evidence

The fact that the structure functions of bound and free nucleons are not equal was discovered in a deep inelastic muon experiment carried out by the European Muon Collaboration at CERN in 1982.

The effect has since then been actively investigated with ever more sophisticated and ingenious deep inelastic experiments with charged leptons. The appeal of such experiments largely consists in the fact that the behavior of the probe – the virtual photon – is well known and calculable.

Other tools have however been exploited, and the subject has been attacked from many angles. In particular attempts have been made to discriminate the roles of sea and valence quarks:

- Neutrinos couple differently to quarks and anti-quarks of a given flavor. Deep inelastic scattering of neutrino and anti-neutrino beams therefore offers a powerful tool to investigate the behavior of sea and valence quarks; unfortunately it is difficult to achieve high statistics with such experiments.
- Inelastic and incoherent production of charmed quarks in DIS mainly proceeds via the photon-gluon fusion mechanism, in which a gluon from the target couples with the virtual photon via a charmed quark. The cross section for this process is proportional to the gluon density in the nucleon.
- Drell-Yan dimuon production is sensitive to the anti-quark distribution in the target nucleon.

It is important to realize that these different techniques may allow to elucidate different facets of the same problem, shedding light on the behavior of quarks and gluons inside nuclear matter.

Before we survey the results so far achieved we shall briefly discuss the only two nuclear effects known prior to the EMC discovery and expected to play a role in deep inelastic scattering: Fermi motion and shadowing.

2.1 Fermi Motion

The nucleons are not stationary in the nucleus, but move with an average momentum k_F .

The structure function F_2 , as we have seen, is proportional in the QPM to the momentum distribution of the quarks inside the nucleus weighted by the square of their charges. This is

only true for a free nucleon, i.e. for a hydrogen target. For a nuclear target the x variable is experimentally determined in the approximation that the nucleon is stationary. Since this is not the case, the structure function measured on nuclear targets is in fact the convolution of the bare nucleon structure function with the momentum distribution function of the nucleon in the nucleus $f_N(z)$:

$$F_2^A(x) = \int_x^A dz f_N(z) F_2^N(x/z), \quad (2.1)$$

where z is the fraction of the nucleus momentum carried by the nucleon multiplied by A .

The effect of Fermi motion on the structure function ratio F_2^A/F_2^D has been calculated by several authors [6]-[10] for various nuclei. Some of the results available at the time of the EMC discovery are shown in fig. 2.1 for F_2^{Fe}/F_2^D . In this figure the solid line is from [6], the dashed line from [7], the dotted one from [8]; the dash-dotted line and triple-dot-dashed line are variants of the calculation [6] and indicate the sensitivity to different model assumptions. Much of the uncertainty follows from the poorly understood high momentum tail of the Fermi distribution.

A detailed discussion of Fermi motion and its interplay with the EMC effect can be found in [10].

2.2 Shadowing

In the low x region ($x < 0.1$), at small Q^2 ($< 1 \text{ GeV}^2$), nuclear effects had long been known to be present and had been actively studied in the 1970's mainly by low energy electron experiments.

The cross section (per nucleon) for interactions between real photons and nuclei decreases with increasing atomic number (cf. [11, 12] and references therein). This effect is known as “shadowing” because of the analogy with the nuclear shadowing of the hadronic cross sections.

As we shall discuss at length below (section 3.2), this can be understood by assuming that the photon fluctuates from its pointlike, “bare photon” state into a superposition of vector mesons like the ρ , the ω or the ϕ , which then interact strongly with the nucleons on the surface of the target nucleons (Vector Meson Dominance, VMD).

A similar but smaller effect had been observed, before the discovery of the EMC effect, for virtual photons in charged lepto-production at small values of Q^2 [13]-[17] (cf. table 2.1 for a summary of the main parameters of the quoted experiments). Figure 2.2 shows some of the data on shadowing in aluminum; the fact that the cross section ratio σ^{Al}/σ^D is lower than unity at small x is ascribed to shadowing of the virtual photon in aluminum nuclei. These early experiments were carried out with low energy electron or muon beams ($< 20 \text{ GeV}$) and covered the small Q^2 region only, typically with $Q^2 < 2 \text{ GeV}^2$. Based on vector meson dominance arguments there was a widespread prejudice that shadowing should disappear at larger Q^2 .

A clear decrease of shadowing with increasing Q^2 – although not as fast as expected – was indeed observed [19] at the Fermi National Accelerator Laboratory (FNAL) near Chicago, using a 209 GeV muon beam, by the only high energy experiment which investigated shadowing before the discovery of the EMC effect. Figure 2.3 shows the results of this experiment in terms of the effective number of nucleons A_{eff} seen by the incident photon:

$$\frac{A_{eff}}{A} = \frac{\sigma^A}{\sigma^D}. \quad (2.2)$$

Lab.-experiment	beam type	beam energy [GeV]	target	x or ν range	Q^2 range [GeV ²]	ref.
FNAL	γ	4-18	C,Cu,Pb	$4 < \nu < 18$ GeV	0	[11]
FNAL	γ	20-185	C,Cu,Pb	$20 < \nu < 185$ GeV	0	[12]
FNAL	γ	90-190	Be,Fe,Pb	$90 < \nu < 190$ GeV	0	[54]
Daresbury	e	5	H,D,Be,C,Al Co,Nb,Ta,U	$2.2 < \nu < 3.8$ GeV	0.066-0.25	[13]
DESY	e	7	H,D,Be,Al,Si	$0.3 < \nu < 6.5$ GeV	0.08-1	[14]
DESY	e	3.08-7	H,C,Al	$\nu < 6.2$ GeV	0.075-1	[17]
Cornell	e		H,D,C,Al Cu,Ta	$2.0 < \nu < 8.5$ GeV	0.1	[16]
SLAC-E61	e	13-20	H,D,Be,Al Cu,Au	$0.02 < x < 0.2$	0.4-1.6	[15]
SLAC-E87	e	9-20	D,Fe	$0.25 < x < 0.9$	3-20	[27]
SLAC-E49B	e	4.5-20	D,Al	$0.075 < x < 0.86$	2-20	[28]
SLAC-E139	e	8-24.5	D,He,Be,C,Al Ca,Fe,Ag,Au	$0.09 < x < 0.9$	2-15	[31, 32]
SLAC-E140	e	4-20	D,Fe,Au	$0.2 < x < 0.5$	1-5	[20]
FNAL	μ	209	C,Cu,Pb	$40 < \nu < 6.2$ GeV	0.01-30	[19]
CERN-EMC NA2	μ	120-280	D,Fe	$0.05 < x < 0.65$	9-170	[5, 26]
CERN-BCDMS NA4	μ	280	D,N	$0.08 < x < 0.7$	26-200	[33]
			D,Fe	$0.2 < x < 0.7$	46-200	
CERN-BCDMS NA4	μ	280	D,Fe	$0.07 < x < 0.65$	14-200	[34]
CERN-EMC NA2'	μ	100-280	D,C,Cu,Sn	$0.03 < x < 0.6$	4-40	[35, 36]
CERN-EMC NA28	μ	280	D,C,Ca	$0.003 < x < 0.1$	0.3-3.2	[37]
FNAL-E665	μ	490	D,Xe	$5 \times 10^{-5} < x < 0.25$	0.01-60	[40, 41]
CERN-NMC NA37	μ	90-200	D,He,Li,C,Ca	$0.0035 < x < 0.65$	0.5-100	[38, 39]
CERN-WA25,WA59	$\nu, \bar{\nu}$	15-160	D,Ne	$0.01 < x < 0.6$	0.05-30	[67, 68]
FNAL-E772	p	800	D,Xe	$0.03 < x < 0.3$	> 16	[72]

Table 2.1: Main parameters of some of the quoted experiments on shadowing and the EMC effect

The results are however somewhat misleading because they are integrated over x . Since large x events are mostly at large Q^2 , the observed Q^2 dependence should in fact be partly ascribed to an x dependence of the signal.

2.3 Experimental Review

The EMC NA2 experiment which discovered that the bound to free nucleon structure function ratio is not what is expected from Fermi motion had been designed to measure absolute structure functions and their Q^2 evolution. The measurement of the iron to deuterium structure function ratio was not included in its experimental program: it was only as a check of the Fermi motion corrections to be applied to the data that such ratio was evaluated.

The result, together with the expected behavior from Fermi motion effects, is shown in fig. 2.4; the theoretical expectations and the data are in clear disagreement. We will comment on this figure more in detail below; we only mention here that the iron and the deuterium data had been collected with different beam energies, intensities, target arrangements and triggers. This is the cause of the rather large systematic errors. Much of the effort of the experiments that followed the EMC result was spent in reducing these errors and extending the kinematic range covered.

A few months after the publication of the EMC ratio, similar results were found at SLAC [27, 28]. They were obtained from the reanalysis of empty-target data from two experiments carried out in the first half of the 1970's to measure the hydrogen and deuterium structure functions. Since the liquid hydrogen and the liquid deuterium had been contained in aluminum and iron vessels, the empty target data – originally taken for background studies – made it possible to measure the aluminum and iron cross sections and to compare them with the deuterium one.

The EMC data were mostly at large Q^2 , while the SLAC ones covered a region of much smaller Q^2 . In spite of the different kinematic conditions, the results were in good agreement with each other in the region of overlap in x . This stimulated the experimental community to further investigate the phenomenon.

Between 1983 and 1985 three experiments were carried out in order to study the EMC effect in different kinematic ranges:

- The E139 experiment at SLAC [31, 32], which used electrons up to 20 GeV and covered the small Q^2 region ($< 15 \text{ GeV}^2$). Cross sections were measured on nine different target materials, from deuterium to gold.
- The NA4 experiment at CERN [33, 34] (BCDMS collaboration), which used 280 GeV muons and covered the very large Q^2 region, from 14 up to 200 GeV^2 . Measurements were carried out on deuterium, nitrogen and iron targets.
- A new experiment by the EMC, NA2' [35, 36], with deuterium, carbon, copper and tin targets, using muon beams with energies from 100 to 280 GeV. The intermediate Q^2 region was covered, $4 < Q^2 < 40 \text{ GeV}^2$.

These experiments were designed with the goal of minimizing the systematic errors in the measurement of cross section ratios: the data on the different materials were taken very close in time, exchanging the targets frequently, or even with several target materials in the beam simultaneously. The experimental conditions under which the data were collected were in all cases kept as similar as possible.

The results of SLAC E139 and BCDMS agreed with the original EMC NA2 ones for x larger than 0.3. There was however some confusion concerning the low x region. Here the EMC NA2 ratios were much higher than the SLAC E139 ones. The BCDMS data were in between and could not solve the issue (see right part of fig. 2.10, below).

The EMC NA2' data clarified the problem (fig. 2.11, below). The structure function ratios at low x were indeed shown to be much lower than those initially measured by NA2 and were smaller than unity for $x < 0.1$.

The agreement with SLAC E139 was at this point satisfactory. This was of deeper relevance than the clarification of an experimental dispute. It confirmed that there was no significant Q^2 dependence of the EMC effect, since the average Q^2 of the EMC NA2' data was much larger than that of the SLAC E139 ones. Furthermore, the fact that the ratio was lower than unity at small x indicated that shadowing persisted at high values of Q^2 .

That shadowing does not depend significantly on Q^2 was confirmed by the EMC NA28 experiment [37], which also used 280 GeV muons but covered the very low x (down to 3×10^{-3}) and Q^2 region (between 0.3 and 3 GeV^2). The experiment had been conceived before the discovery of the EMC effect, and was intended to carry out a systematic investigation of shadowing.

The combination of these data with the EMC NA2' and the SLAC E139 ones at higher x provided finally a consistent experimental picture of nuclear effects in DIS, in which shadowing and the EMC effect join smoothly in the region of $x \approx 0.2$ and there is little dependence on Q^2 .

About at the same time another SLAC electron experiment, E140 [20], measured the difference of $R = \sigma_L/\sigma_T$ in e -D, e -Fe, and e -Au scattering. The results were consistent with no dependence of R on A , lending experimental support to the identification of cross section and structure function ratios.

Recently two experiments have added valuable information in the low x region, NMC NA37 at CERN and E665 at FNAL.

- NA37 [38], carried out by the New Muon Collaboration (NMC), measured the x and Q^2 dependence of the ratios He/D, C/D and Ca/D in the range $0.0035 < x < 0.65$ and $0.5 < Q^2 < 90 \text{ GeV}^2$.

The same collaboration also measured the ratios C/Li, Ca/Li and Ca/C in the range $0.0085 < x < 0.6$ and $0.8 < Q^2 < 17 \text{ GeV}^2$ [39].

- E665 [40, 41] measured the Xe to D cross section ratio down to $x = 5 \times 10^{-5}$ and $Q^2 < 0.01$.

Means other than the study of charged lepton DIS have been used to investigate the effects of the nuclear environment on the nucleon:

- Inelastic and incoherent production of charm-anti-charm quark pairs in deep inelastic scattering, as well as in real photo-production experiments, is thought to proceed mainly via the fusion of the photon with a gluon from the target nucleon, with the photon and the gluon both coupling to the same charmed quark. In this framework the cross section for J/ψ production is proportional to the gluon distribution function in the nucleon. The difference between the gluon distributions in bound and free nucleons can thus be inferred from the ratio of the respective cross sections for J/ψ production.

The experimental situation in the field was until recently somewhat confused, because of apparently contradictory results from different experiments. New results from NMC NA37 at CERN [56] have largely clarified the matter and can be interpreted in terms of an enhancement of the gluon distribution in tin with respect to that in carbon for $0.02 < x < 0.2$. Unfortunately the statistical errors associated with J/ψ production data are still very large if compared with those from the standard structure function analyses.

- Experiments with neutrino beams are able to discriminate between sea and valence quarks, and are therefore potentially very interesting. Unfortunately such experiments are affected by large uncertainties and so far no single experiment could by itself confirm the effect. Just to mention one problem, while several hundred thousand events after cuts are not uncommon for a charged lepton scattering experiment, the largest number of events accumulated by a neutrino experiment is about 40,000.
- Drell-Yan experiments are sensitive to the anti-quark distribution in the nucleon. In one of them [72] shadowing effects similar to those seen in DIS were observed.

We shall now examine the available experimental results in more detail. We shall first present the recent evidence on the lack of nuclear effects in the ratio $R = \sigma_L/\sigma_T$; we will then concentrate on the effects on F_2 , following an approximately chronological order. The main parameters of the experiments described are summarized in table 2.1.

2.3.1 Nuclear Effects on $R = \sigma_L/\sigma_T$

Some experiments measured cross section ratios rather than F_2 ratios. The two ratios are the same only as long as $R = \sigma_L/\sigma_T$ is the same for different target materials. This can be seen by rewriting relation 1.4 for a nuclear target A and for D, and then taking the ratio.

The SLAC experiment E140 published in 1988 results of a measurement of R on deuterium, iron and gold targets in the kinematic range $0.2 < x < 0.5$ and $1 < Q^2 < 5 \text{ GeV}^2$ [20, 21], carried out at the SLAC electron beam with energies between 4 and 20 GeV. While providing the first clear evidence for a fall-off of R with increasing Q^2 [22, 23], the results demonstrate that $R^A - R^D$ is consistent with zero; the average value of $R^A - R^D$, for A=Fe and Au is $\langle R^A - R^D \rangle = 0.001 \pm 0.018 \text{ (stat.)} \pm 0.016 \text{ (syst.)}$. Figure 2.5 shows $R^A - R^D$ as a function of x . From these data one can also deduce $\langle R^{Au} - R^{Fe} \rangle = 0.020 \pm 0.044$, where statistical and systematic errors are summed in quadrature.

A similar conclusion was recently obtained by the NMC [24], based on their Ca/C data at 200 [38] and 90 GeV [39]. Their preliminary result is that $\langle R^{Ca} - R^C \rangle = 0.027 \pm 0.026 \text{ (stat.)} \pm 0.020 \text{ (syst.)}$ at $\langle Q^2 \rangle = 6 \text{ GeV}^2$, for $0.007 < x < 0.2$.

Nuclear effects seem therefore to affect the structure functions F_2 and F_1 in the same way and to cancel in R .

2.3.2 The Early Results: EMC NA2 and the SLAC Reanalysis

The ratio of the iron to the deuterium structure function was originally evaluated by the EMC (NA2) [5] as a direct check of the procedure used to correct for Fermi motion the structure functions obtained from heavy targets. The results were thus expected to resemble the curves shown in fig. 2.1.

For iron, the F_2 data obtained with incident muon energies of 120, 200, 250 and 280 GeV [25] were used. For deuterium, data at 280 GeV only were available [26]. Note that the Fe and D measurements were taken under different experimental conditions and at different times.

The iron data were corrected for the non-isoscalarity of $^{56}_{26}\text{Fe}$; neither of the data sets was corrected for Fermi motion. Only data points with a total systematic error of less than 15% were used for the ratio.

The kinematic region covered is limited by the extent of the deuterium data. The x range extends from $x = 0.05$ to $x = 0.65$. The Q^2 range is different for each x value and varies from 9 to 27 GeV^2 for $x = 0.05$, from 11.5 to 90 GeV^2 for $x = 0.25$ and from 36 to 170 GeV^2 for $x = 0.65$.

The results are presented in fig. 2.4 as a function of x , together with the behavior expected from Fermi motion calculations. The disagreement between experiment and theoretical expectations is quite striking. No significant Q^2 dependence was observed.

The EMC measurement was subsequently confirmed by a reanalysis of data taken ten years earlier by the SLAC experiments E87 [27] and E49B [28], a remarkable example of archaeology in high energy physics. The experiments had been designed to extract the proton and the

neutron structure functions; their results had been initially published in 1973 for E49B [29] and in 1974 for E87 [30]. The data had been collected using electron beams with energies ranging between 9 and 20 GeV for E87, 4.5 and 20 GeV for E49B.

After the EMC discovery the deuterium events were compared with those originating in the deuterium vessel walls, made of steel in the case of E87 and of aluminum in that of E49B. The kinematic range covered is $0.25 < x < 0.9$, $3 < Q^2 < 20 \text{ GeV}^2$ for E87 and $0.075 < x < 0.86$, $2 < Q^2 < 20 \text{ GeV}^2$ for E49B. The resulting cross section ratios are shown in fig. 2.6 (full symbols), together with the EMC data and the photo-production results [11]. The results of the SLAC experiment E61 [15] (labelled "SLAC") on shadowing are also shown; this experiment was carried out in the 1970's with electrons up to 20 GeV and covered the range $0.1 \text{ GeV}^2 < Q^2 < 1.8 \text{ GeV}^2$ (targets: H, D, Be, Al, Cu, Au).

2.3.3 The Second Generation of Experiments: SLAC E139, BCDMS and EMC NA2'

The first experiment specifically designed to study the EMC effect was SLAC E139 [31], which took data in the range $0.09 < x < 0.9$ and $2 < Q^2 < 15 \text{ GeV}^2$ on deuterium and He, Be, C, Al, Ca, Fe, Ag, Au targets. Electron beams with energies ranging from 8 to 24.5 GeV were used in conjunction with the SLAC 8 GeV spectrometer. The experimental set-up did not allow simultaneous measurements with several targets, but the deuterium and the heavy targets were frequently exchanged in order to minimize systematic errors.

The cross section ratios σ^A/σ^D found by this experiment are presented in fig. 2.7. The figure also shows the original EMC NA2 results [5], those of the SLAC E49B [28] and E87 [27] experiments (labelled "Bodek *et al.*"), just described, as well as those of two experiments on shadowing: SLAC E61 [15] (labelled "Stein *et al.*"), mentioned above, and the FNAL muon experiment [19] also already quoted before (labelled "Goodman *et al.*").

While the original EMC NA2 results are again confirmed for $x > 0.3$, the two data sets rather clearly disagree at lower values of x , the EMC points being systematically higher – by up to 15% – than the SLAC ones.

In the E139 results Fermi smearing is apparent at large values of x . These are in fact the only data in which the Fermi motion behavior is clearly visible.

The E139 data have recently been the subject of a reanalysis [32]; the results of the reanalysis are shown in fig. 2.8, together with the old ones. The main differences are in the radiative corrections (this affects mainly the small x region), in the statistics (more data have been included) and in the analysis procedure itself. The results differ at small x , where the new points are higher than the old ones by up to a few per cent, and in some cases (e.g. Ca) also at larger values of x , with the new results being again higher than the old ones.

Similar changes with respect to the old E139 results had already been found by the SLAC experiment E140 [20], designed to investigate the A dependence of $R = \sigma_L/\sigma_T$ and discussed above; in the same paper the results of an earlier reanalysis of the E139 data for Fe/D using an improved radiative corrections procedure had also been presented.

At CERN the EMC results were later confirmed by the BCDMS collaboration (experiment NA4), which also used a 280 GeV muon beam, the same as that of EMC. The apparatus consisted of an instrumented iron toroid with large acceptance for high Q^2 events. Data were taken with two target materials simultaneously in the beam, thus minimizing systematic errors from uncertainties on the energy calibration of incoming and scattered muons, on the detector efficiencies and on the beam flux measurements.

Two sets of measurements were carried out. In the first one [33] the target pairs simultaneously in the beam were Fe/D and Fe/N. The ratios of the nitrogen and iron to the deuterium structure functions were measured in the range $0.08 < x < 0.7$ and $0.2 < x < 0.7$, respectively. The Q^2 range covered is 26 to 200 GeV² for the F_2^N/F_2^D data and 46 to 200 GeV² for the F_2^{Fe}/F_2^D ones. The results are shown in fig. 2.9 together with the original EMC NA2 ones [5] for F_2^{Fe}/F_2^D and those, just presented, from SLAC E139 [31].

The same collaboration later published a high statistics study of F_2^{Fe}/F_2^D over a wider x and Q^2 region [34]. The range covered was $0.07 < x < 0.65$ and $14 < Q^2 < 200$ GeV². The results are shown in fig. 2.10 and are compared with the earlier BCDMS ones [33] (open circles) and with those of the EMC NA2 [5], SLAC E139 [31] (labelled “Arnold *et al.*”) and SLAC E61 [15] experiments (labelled “Stein *et al.*”).

The BCDMS results were explicitly obtained as ratios of structure functions: the measured yields were first converted to cross sections, after correcting for acceptance, detector efficiencies etc., and then to structure functions.

The BCDMS data agree with the EMC NA2 and SLAC E139 ones at large x . For $x < 0.2$ they indicate a ratio lower than that originally measured by EMC but still above the E139 points, thus leaving open the question of the behavior of the EMC effect at small x . No Q^2 dependence is visible over the measured range.

The EMC collaboration themselves engaged in 1984 and 1985 in two series of new measurements within the NA2' experiment.

For the first measurement [35], data were taken on deuterium, carbon, copper and tin targets. The geometry of the targets was similar and they were frequently interchanged under the same beam and spectrometer conditions. Spectrometer acceptance corrections were the same for all targets and thus canceled in the ratio. Also, the effects of time dependent inefficiencies in the spectrometer and in the beam intensity monitoring apparatus could to a large extent cancel out.

Data were taken at different incident muon energies, ranging from 100 to 280 GeV. The kinematic range covered is $0.03 < x < 0.6$ and $4 < Q^2 < 40$ GeV². The results are shown in fig. 2.11 with full circles.

For the second measurement [36], data were taken with deuterium and copper targets simultaneously in the beam. Both targets were therefore exposed to the same beam flux and the ratios were thus free of normalization uncertainties due to the beam flux measurement. Acceptance corrections were instead different for each target, since the acceptance varied as a function of the target position. The kinematic range is similar to that covered by the first measurement. The results are shown in fig. 2.11 with open circles; the merged results for the Cu/D ratio are shown in fig. 2.12; this is, at the moment, the most precise measurement of nuclear effects in copper.

The EMC NA2' data agree with the previous high Q^2 results (the original EMC NA2 one [5], and the BCDMS ones [33, 34]), at large x . The new results indicate however a turning over of the ratio at low x , similar to that observed by the lower energy experiment SLAC E61, mentioned above, thus providing the first clear evidence that shadowing persists also at considerable values of Q^2 . This was largely unexpected because shadowing was believed – according to VMD – to disappear quickly for $Q^2 > 1$ GeV².

2.3.4 The Shadowing Region: the EMC NA28 Results

A rather clear confirmation of the fact that shadowing depends only weakly – if at all – on Q^2 came from yet another EMC experiment: NA28.

In this experiment a special low angle trigger system was used, which was capable of selecting events where the muon was scattered at very small angles, down to 2 mrad. The incident muon energy was 280 GeV. Data were taken with deuterium, carbon and calcium targets. Structure functions were extracted for each target individually and then ratios were taken.

The ratios of the carbon and calcium to deuterium structure functions were measured in the range $0.003 < x < 0.1$ and $0.3 < Q^2 < 3.2 \text{ GeV}^2$ [37]. The results are presented in fig. 2.13 as a function of x . No Q^2 dependence is visible in the data.

The large systematic errors are a consequence of the uncertainty in the trigger efficiency and of the fact that the targets were not simultaneously in the beam, nor at the same position: both the uncertainty on the flux and that on the acceptance had thus to be taken into account.

The figure also shows the results of other lepto-production experiments discussed before: EMC NA2' [35] (labelled “Ashman”), BCDMS [33, 34] (labelled “Bari” and “Benvenuti”), SLAC E61 (labelled “Stein”) and the FNAL muon experiment [19] (labelled “Goodman”). For the latter the x values of the points were computed by dividing the published Q^2 values by $2M\nu$ using $\nu = 150 \text{ GeV}$, an average value in the acceptance region of the experiment ($110 < \nu < 200 \text{ GeV}$ for $Q^2 < 3 \text{ GeV}^2$ and $40 < \nu < 200 \text{ GeV}$ for $Q^2 > 3 \text{ GeV}^2$). The points on the vertical axis are ratios from the real photo-production experiment quoted above [12], where the energy of the photon beam was 60 GeV.

2.3.5 Recent Developments: NMC and E665

NMC NA37

The New Muon Collaboration at CERN carried out, in the period 1986-1989, a number of experiments investigating the structure functions of free and bound nucleons.

The experimental set-up was a modified and upgraded version of the EMC spectrometer. The modifications were meant to improve the ability of the spectrometer to measure both absolute cross sections and ratios of cross sections. For the latter, the special target arrangement was instrumental.

A complementary target set-up was used, made up of two target sets. In the simplest case each set consisted of two targets of different materials, simultaneously exposed to the beam, one behind the other along the beam direction. This set was frequently exchanged with a complementary one where the two target materials were interchanged. Geometrical acceptance and efficiency corrections therefore cancel in the calculation of the ratios as do the beam fluxes. Furthermore the frequent exchange of the two target sets (every 30 min, a total of a few hundred times for each data set) substantially eliminated the effects of possible time dependences of the apparatus acceptance and efficiency. Systematic errors due both to the beam flux measurement and to acceptance corrections – two of the most important sources of systematic uncertainties in previous experiments – were thus reduced to essentially zero.

The sensitivity of the trigger system was extended with respect to that of the EMC apparatus. Two independent trigger systems were in fact used. Muon scattering angles larger than 12 mrad were covered by the large angle trigger, essentially the same as that of EMC. A small-angle trigger was sensitive to muons scattered between 5 and 15 mrad. Both triggers

covered the region $x < 0.4$, the small angle trigger events having smaller Q^2 and ν . Events at larger x were accepted by the large angle trigger only.

The first set of results concerns the measurement of the structure function ratios F_2^A/F_2^D for the isoscalar nuclei ${}^4\text{He}$, C and Ca, at an incident muon energy of 200 GeV [38]. The kinematic range $0.0035 < x < 0.65$ and $0.5 < Q^2 < 90 \text{ GeV}^2$ is explored. The Q^2 range covered varies with x : for example at $x = 0.0035$ $\langle Q^2 \rangle \approx 0.7 \text{ GeV}^2$, at $x = 0.01$ $\langle Q^2 \rangle \approx 2 \text{ GeV}^2$, at $x = 0.1$ $\langle Q^2 \rangle \approx 10 \text{ GeV}^2$ and at $x = 0.2$ $\langle Q^2 \rangle \approx 20 \text{ GeV}^2$.

Figure 2.14 shows the results. Both the statistical and the systematic errors at small x are typically 1%.

The results for He/D are the first to cover the low x region. In the same region the C/D and Ca/D ratios were previously measured by the EMC NA28 [37], but the NMC data are significantly more precise.

In the intermediate x region, where the data were so far inconclusive, an enhancement of the C/D and Ca/D ratios above unity is now clearly visible. The data suggest also some A dependence of the enhancement size.

At large x , the NMC results are consistent with the well known EMC effect shape.

In the measured Q^2 range there is little indication of a Q^2 dependence of the ratios. The logarithmic slopes b obtained from fits of the type $F_2^A/F_2^D = a + b \ln Q^2$ were evaluated in each bin separately. The slopes are plotted in fig. 2.15, together with those extracted from the recent reanalysis of the SLAC E139 data [32]. For He/D and C/D no significant Q^2 dependence is observed. In the case of Ca/D negative slopes are found over a large fraction of the x range. Only in a few bins around $x = 0.07$ -0.1, however, are they different from zero in a statistically significant way. This is marginally in agreement with the behavior observed by SLAC E139 for Fe/D (squares in the Ca/D plot) and for Au/D, also shown in fig. 2.15.

In a different experiment, the NMC measured the structure function ratios F_2^C/F_2^{Li} , F_2^{Ca}/F_2^{Li} and F_2^{Ca}/F_2^C at an incident muon energy of 90 GeV, covering the kinematic range $0.0085 < x < 0.6$ and $0.8 < Q^2 < 17 \text{ GeV}^2$ [39]. This combination of targets allows the comparison of structure functions for pairs of isoscalar nuclei differing either in radius or in density. In particular ${}^6\text{Li}$ ($r = 2.6 \text{ fm}$, $\rho = 0.04 \text{ fm}^{-3}$) and ${}^{12}\text{C}$ ($r = 2.5 \text{ fm}$, $\rho = 0.09 \text{ fm}^{-3}$) have equal radii but different densities, whereas ${}^{12}\text{C}$ and ${}^{40}\text{Ca}$ ($r = 3.5 \text{ fm}$, $\rho = 0.11 \text{ fm}^{-3}$) differ more in radius than in density.

Figure 2.16 shows the results for the F_2^C/F_2^{Li} , F_2^{Ca}/F_2^{Li} and F_2^{Ca}/F_2^C , while fig. 2.17 shows the F_2^{Li}/F_2^D obtained by comparing these results with those of fig. 2.14. The following observations can be made from the data.

(i) At low x there is a depletion of the structure function in C relative to that in Li, which has the same radius but a lower mean density. However the structure function depletion in Ca relative to that in Li is twice as large, implying that the depletion increases with both radius and density.

(ii) At intermediate x , the structure functions in the nuclei with high density (C and Ca), exhibit similar enhancements of 1-2% relative to that in Li. There are no significant differences between the C and Ca structure functions for x between 0.07 and 0.4. This indicates a weak dependence of the structure function on the density with little or no sensitivity to the radius in the intermediate x region.

(iii) The Li/D ratio is remarkably similar to the He/D one (also shown in fig. 2.17). For small x this observation is in agreement with the conclusion reached in (i): Li as compared to He ($r = 1.7 \text{ fm}$, $\rho = 0.09 \text{ fm}^{-3}$) has a much larger radius but a much smaller density, so that

in this region two opposing dependences should tend to cancel. In the intermediate x region the weakness of the radius and density dependence precludes, in view of the experimental uncertainties, a useful comparison with the conclusion suggested in (ii).

FNAL E665

Preliminary results from the muon scattering experiment E665 at the FNAL Tevatron on the ratio of xenon ($A = 131$) to deuterium cross sections were recently presented [40]. The experimental apparatus is an open geometry muon spectrometer with acceptance for forward tracks down to zero scattering angle. The data cover the range $0.001 < x < 0.25$ and $0.1 < Q^2 < 20 \text{ GeV}^2$; they were taken at an incident muon energy of 490 GeV. The results are shown in fig. 2.18 as a function of x and Q^2 . No significant Q^2 dependence is visible in these data either and the observed x dependence agrees with that measured by previous experiments.

Figure 2.19 shows the results of a different analysis of the E665 data [41], in which the x and Q^2 ranges were extended to lower values: $5 \times 10^{-5} < x < 0.2$ and $0.01 < Q^2 < 60 \text{ GeV}^2$. Radiative events in which the muon radiates a photon and elastic μ - e scattering events were rejected by using the information of a Pb-MWPC electromagnetic calorimeter. The remarkable feature of these data is that the ratio of the xenon to the deuterium cross section appears to saturate for x values lower than $\approx 10^{-3}$. The Q^2 values where saturation is observed are very small, and range from $Q^2 = 0.01 \text{ GeV}^2$ to about 0.7 GeV^2 . The very use of the Bjorken variable x at such low values of Q^2 has been questioned [137].

The results were confirmed by an independent analysis in which only events with at least two positive or two negative tracks (in addition to the scattered muon) fitted to the interaction vertex were retained. Such cut explicitly excludes μ - e elastic events and e^+e^- production in the target; it also implicitly assumes that high multiplicity events are due to inelastic scattering.

2.3.6 The A dependence of F_2^A/F_2^D

To a first approximation, the dependence of F_2^A/F_2^D on the atomic number A can be described by a function of the type $F_2^A/F_2^D = \alpha + \beta \ln A$ in each x bin separately (cf. e.g. [42]).

Figure 2.20 shows the results of the fits for the slopes $\beta = d(F_2^A/F_2^D)/d \ln A$ as a function of x , obtained by using the NMC data [38] on He/D, C/D and Ca/D and the SLAC E139 ones [32] on He/D, Be/D, C/D, Al/D, Ca/D, Fe/D, Ag/D and Au/D. The NMC and SLAC data have little overlap; together they provide an accurate picture of the A dependence of nuclear effects in DIS.

The figure shows that at small x the size of shadowing in each x bin increases with $\ln A$. For $0.07 < x < 0.3$, approximately within the enhancement region, the ratios are A independent or slightly increasing with A . For $0.2 < x < 0.8$ the ratio decreases with $\ln A$.

2.3.7 The Integrals of the Structure Function Differences $F_2^A - F_2^D$

From the structure function ratio and the absolute structure function F_2^D , the difference $F_2^A - F_2^D$ can be evaluated as $(F_2^A/F_2^D - 1)F_2^D$ and thus the integral $\int (F_2^A - F_2^D)dx$ can be computed. In the framework of the quark-parton model this integral, evaluated over the full x range, represents the difference of the momentum fraction carried by charged partons in bound nucleons relative to that for deuterium.

We present here the integrals computed by combining the NMC results for He/D, C/D and Ca/D [38] with those of the reanalysis of the SLAC E139 data for the same nuclei [32].

	x range	$\int \left(\frac{F_2^A}{F_2^D}(x) - 1 + f_M \right) F_2^D(x) K_2(x) dx$	stat.	syst.
He/D	0.0035 – 0.25	0.6×10^{-3}	0.3×10^{-3}	0.6×10^{-3}
	0.0035 – 0.80	-0.5×10^{-3}	0.4×10^{-3}	2.2×10^{-3}
C/D	0.0035 – 0.25	0.7×10^{-3}	0.3×10^{-3}	0.5×10^{-3}
	0.0035 – 0.80	-0.7×10^{-3}	0.4×10^{-3}	1.3×10^{-3}
Ca/D	0.0035 – 0.25	0.0×10^{-3}	0.3×10^{-3}	0.6×10^{-3}
	0.0035 – 0.80	-1.5×10^{-3}	0.4×10^{-3}	1.4×10^{-3}

Table 2.2: Integrals of the structure function differences $F_2^A - F_2^D$ at $Q^2 = 5 \text{ GeV}^2$. For the region $0.0035 < x < 0.25$ the NMC results [38] only are used; for larger values of x the SLAC E139 results [32] only are used

The evaluation of this integral involves some subtleties. For a detailed discussion of the method we refer the reader for instance to [42]. Here we only recall that the quantity in which we are interested is actually

$$M_2^A - M_2^D = \int_0^2 \left(\frac{F_2^A}{F_2^D} - 1 + f_M \right) F_2^D K_2(x, Q^2) dx. \quad (2.3)$$

In this formula the constant $f_M = (AM/M_A - 2M/M_D)$ is due to the mass defect of bound nucleons (M_A , M_D are the masses of nucleus A and of deuterium, respectively). In order to account for target mass effects the integration must be carried out over the Nachtmann scaling variable ξ^A [43, 44]. The function $K_2(x, Q^2)$, close to unity, includes the jacobian of the transformation $\xi^A \rightarrow x$.

In order to evaluate the integrals, the experimental results for the structure function ratios were used, while a parametrization was used for F_2^D , based on a fit [45, 42] to the results of deep inelastic scattering experiments [37, 4, 46, 47] and to low energy data in the resonance region [48]. The parametrization was evaluated at $Q^2 = 5 \text{ GeV}^2$; this value is very close to the average Q^2 of both the NMC and the SLAC E139 data. No Q^2 dependence was assumed for the ratios.

Table 2.2 shows the results; they are based on the NMC data in the region $0.0035 < x < 0.25$ and on the reanalyzed SLAC E139 ones for x up to 0.8. The systematic errors include the contributions of the systematic errors on the measured points and their normalization uncertainty, as well as the uncertainty on F_2^D [45, 42].

The results are different from those presented in [38], where the old SLAC E139 data [31] were used. The effect of including the reanalyzed SLAC points is to bring the values of the integrals closer to zero.

Note that the available data do not allow the evaluation of the integrals over the full x range. In the range covered by the NMC and SLAC results the integrals are compatible with zero.

2.3.8 Direct Measurements of the Gluon Distributions in Nuclei: J/ψ Production

An entirely different way to investigate nuclear effects in deep inelastic scattering of charged leptons is to probe directly the gluon distribution of different targets. This can be done by studying the inelastic production of J/ψ mesons.

Inelastic charm-anti-charm production is assumed to proceed according to the photon-gluon fusion mechanism [49] in which the virtual photon couples to a gluon in the target via a charm-anti-charm quark pair. A specific formulation of the photon-gluon fusion idea is the Color Singlet (CS) model (fig. 2.21) [50, 51, 52], in which the J/ψ is represented by a definite wave function of the $c\bar{c}$ system, describing a color singlet state with $J^P = 1^-$ and the J/ψ rest mass. In this model the amplitudes of the process $\gamma g_1 \rightarrow J/\psi g_2$ (fig. 2.21) are calculated under the assumption that the virtual photon fuses with a hard gluon g_1 from the target nucleon; in the calculation color conservation is guaranteed by the emission of the (hard) gluon g_2 . In this framework the inelastic J/ψ production cross section is thus directly proportional to the gluon density in the nucleon. Hence the ratio of the gluon densities for different nuclei can be obtained from the ratio of the respective J/ψ production cross sections.

In addition to Q^2 and ν , two other variables are used to describe J/ψ production: $z = E_{J/\psi}/\nu$ is the fraction of the virtual photon energy carried by the J/ψ in the laboratory frame; p_T^2 is the transverse momentum squared of the J/ψ with respect to the photon direction.

The EMC and the E691 Results

A short time after the discovery of the EMC effect, the EMC themselves measured the total cross section for inclusive J/ψ photo-production (extrapolated to $Q^2=0$) in hydrogen, deuterium (at 280 GeV incident muon energy) and iron targets (at 250 GeV incident muon energy) [53]. A correction was applied to the Fe data at high z ($z > 0.95$) and small p_T^2 ($p_T^2 < 0.18 \text{ GeV}^2$) in order to take into account the fact that in this region a fraction of the observed J/ψ 's are produced coherently.

The ratio of iron to hydrogen and deuterium cross sections was found to be larger than unity by 1.44 ± 0.12 (stat.) ± 0.20 (syst.), which was interpreted as an enhancement of the gluon distribution in iron with respect to that in H and D.

A year later however, the experiment E691 [54], carried out at the FNAL photon beam (with $E_\gamma = 80\text{-}190 \text{ GeV}$), produced results in contradiction with those of the EMC and rather indicating a depletion of the photo-production cross section on iron compared to that on berillium by a factor 0.79 ± 0.08 . In this experiment coherent effects were eliminated by removing events at small p_T^2 , $p_T^2 < 0.15 \text{ GeV}^2$.

The discrepancy remained unresolved until the appearance of the recent results [55, 56, 57] from the NMC experiment at CERN, which we now present in some detail.

The NMC Result

The NMC studied J/ψ production on hydrogen, deuterium, carbon and tin targets using a high energy muon beam (200 and 280 GeV). The J/ψ was identified by its decay into a pair of oppositely charged muons; the momenta of these muons, and that of the scattered muon, were measured for each event. The H and D targets, as well as the C and Sn ones, were arranged in a complementary set-up similar to that described above for the F_2 measurement, allowing cancellation of beam fluxes and acceptance corrections in the evaluation of cross section ratios.

Inelastically produced J/ψ 's were selected by requiring z to be less than 0.90. The cross sections measured on hydrogen and deuterium targets were found [55] to be in good agreement with each other and to be reproduced by the CS model, except for an overall normalization factor.

Having established that the CS model offers a good description of the data on hydrogen and deuterium – at least in shape – and therefore that the observed J/ψ yield is indeed proportional to the gluon density in the nucleon, the influence of the nuclear medium on the gluon distribution was then studied by comparing the inelastic cross sections measured on tin and carbon [56]. The ratio of the gluon distributions $R_{in}(\text{Sn/C}) = G_{\text{Sn}}(x)/G_{\text{C}}(x)$ thus found is shown as a function of x in fig. 2.22. The average value of $R_{in}(\text{Sn/C})$ over the range $0.02 < x < 0.2$ is 1.13 ± 0.08 , indicating an enhancement of the gluon distribution in tin with respect to that in carbon.

The Resolution of the Discrepancy

The NMC result is apparently in contradiction with the previous ones from EMC and E691, which in turn do not agree with each other. The NMC themselves offered a possible solution [56] of these discrepancies, also shedding light on the disagreement between the EMC and E691 results.

Production of J/ψ 's in muon interactions with nuclei includes several contributions: the production cross section can be considered as the sum of the cross sections for coherent, quasi-elastic and inelastic interactions. In coherent J/ψ production (at large z and small p_T^2) the recoiling nucleus emerges intact from the interaction, whereas in quasi-elastic production (at large z and large p_T^2) the interaction occurs elastically with a nucleon in the nucleus. Since the mechanisms are different, one may expect the corresponding cross section ratios for heavy and light nuclei to be different.

Indeed NMC measured [56] $R_{coh}(\text{Sn/C}) = 1.54 \pm 0.07$ for coherent J/ψ production ($z > 0.9$ and $p_T^2 < 0.3 \text{ GeV}^2$), $R_{qe}(\text{Sn/C}) = 0.79 \pm 0.06$ for quasi-elastic production ($z > 0.9$ and $p_T^2 > 0.4 \text{ GeV}^2$) and, as we said already above, $R_{in}(\text{Sn/C}) = 1.13 \pm 0.08$ for inelastic production, which is the dominant process only when z is small ($z < 0.85$) and p_T^2 is large ($p_T^2 > 0.4 \text{ GeV}^2$). It is then essential in the analysis to be able to separate the kinematic regions where each of the above mentioned mechanisms is important. In [56] it is argued that, due to a lack of resolution in z and p_T^2 , the previous experiments could not completely isolate the inelastic contribution, their J/ψ samples thus being contaminated by a sizable quasi-elastic or coherent contribution.

In particular in the EMC experiment the resolution in z and p_T^2 did not allow a full subtraction of the coherent contribution; as a consequence coherent events from iron were included in the inelastic sample, biasing the ratio of iron to hydrogen and deuterium cross sections towards a higher value.

Likewise in the E691 experiment the energy fraction z was not measured and only coherent effects were removed by excluding events with $p_T^2 < 0.15 \text{ GeV}^2$; as a result the measured yield of J/ψ events contained contributions both from the inelastic and the quasi-elastic processes. The value of the ratio R_{qe} significantly below unity measured by the NMC allows one to understand the depletion effect observed in the FNAL experiment.

Since the various experiments investigated different target pairs, these conclusions are valid only if the A dependence in the measured ratios is small compared to the effects of going from one kinematic region (e.g. coherent, quasi-elastic etc.) to another.

The NMC results can be compared with those on the production of J/ψ and Υ mesons in

nuclear targets as observed for proton and pion beams of energies from 200 to 800 GeV [59]-[61]. Under some assumptions [58] the nuclear effects seen there can be ascribed to modifications of the gluon distribution. Figure 2.23 compares such results with the NMC ones [56]. The NMC and hadron beam results are complementary in their kinematic coverage. Figure 2.24 shows similar results for the A/D and A/H ratio.

2.3.9 Neutrino Scattering Results

Deep inelastic neutrino scattering is in principle an ideal tool to investigate the EMC effect as it may allow the behavior of the sea and valence quarks to be separated experimentally. The statistical significance of neutrino experiments is however poor and no individual experiment [62]-[67] could so far, by itself, confirm the effect.

In fig. 2.25 we show the results of a recent high statistics experiment carried out at the CERN bubble chamber BEBC by the WA25 and WA59 collaborations [67]. It represents the state of the art in the measurement of nuclear effects in neutrino DIS interactions. The ratios of the neon ($A = 20$) to deuterium cross sections are shown as a function of x and y for the neutrino, the anti-neutrino and the combined data samples. The results of the SLAC electron experiment E139 [31] for σ^{Al}/σ^D are also shown.

The following observations can be made on these data:

1. The ratio of the x distributions for $0.2 < x < 0.6$ is similar to that measured in charged lepton DIS.
2. The ratio of the y distributions measures the difference in the sea quark distribution relative to the valence one, averaged over x . The data indicate that the shape of the sea quark distribution is the same for Ne and D; they suggest however that the normalization of the distribution is smaller in Ne than in D by about $10\text{-}15\% \pm 5\text{-}10\%$.
3. For $Q^2 > 1 \text{ GeV}^2$ no Q^2 dependence is visible in the measured ratios.
4. Finally, the total neutrino cross sections on neon and on deuterium turn out to be the same at the 4% level. This indicates that, within 4%, the integral gluon momentum fraction, as well as the integral quark momentum fraction, is the same in neon and deuterium.

The same experiment also investigated shadowing effects in neutrino and anti-neutrino interactions on neon and deuterium [68].

Similarly to virtual photons, $W^\pm$ bosons may fluctuate into a virtual hadron system composed of vector and axial-vector mesons. Partial Conservation of Axial Current (PCAC) [69] also predicts a pion contribution to the propagator of the weak current; this contribution is important for $Q^2 < 0.2 \text{ GeV}^2$. For larger values of Q^2 , the vector (ρ) and axial-vector (a_1) contributions become dominant.

Figure 2.26 presents the measured ratio of the neon to deuterium cross sections for $x < 0.2$ as a function of Q^2 for the neutrino, the anti-neutrino and the combined data samples. Figure 2.27 shows the x dependence of the ratio for the combined neutrino and anti-neutrino data in different Q^2 bins. The full curves represent the predictions derived from PCAC. It is clear that the cross section in neon is indeed depleted relative to that in deuterium at low Q^2 . This is the first evidence of shadowing in neutrino interactions.

Recently a new method to investigate nuclear effects in neutrino interactions was proposed [70]. High energy neutrino interactions in a heavy liquid bubble chamber can be classified by the presence or absence of dark tracks at the interaction vertex. Dark tracks are presumed to be mostly due to slow protons emerging from the target nucleus and may indicate an interaction with a deeply bound nucleon, rather than with a loosely bound one on the surface of the nucleus. Quark distributions can be extracted from the two sets of data and the results compared; those with dark tracks show an enhanced nuclear effect. The method is however not uncontested and there are claims that the effects seen are not correlated to nuclear phenomena and that they may be detected with hydrogen and deuterium targets as well [71].

2.3.10 Nuclear Effects in Drell-Yan Dimuon Production

Continuum dimuon production in high-energy proton-nucleus collisions provides an independent measurement of the modification of the quark distribution in nuclei. For fractional longitudinal momentum $x_F > 0.2$, the dominating mechanism in proton induced Drell-Yan production is the quark-anti-quark annihilation process $q_{\text{proton}} + \bar{q}_{\text{target}} \rightarrow \mu^+ \mu^-$. The dimuon yield is thus proportional to the anti-quark content of the target nucleon. Some caution should be however exercised in comparing these results to the DIS ones, as it is not firmly established that the shadowing term in DIS can be incorporated in the standard factorization theorems relating DIS and Drell-Yan [74, 137].

The FNAL experiment E772 [72, 73] observed a depletion of the dimuon yield in calcium, iron and tungsten targets with respect to deuterium for $0.03 < x < 0.3$; the Q^2 range excludes the quarkonia resonances: $16 < Q^2 < 81 \text{ GeV}^2$ and $Q^2 > 121 \text{ GeV}^2$. Note however that the Q^2 scales in Drell-Yan and DIS are not a priori necessarily comparable. Figure 2.28 shows the results; the W/D data are compared with the EMC NA2' ones for Sn/D [35]. Figure 2.29 shows the comparison with the Ca/D ratio measured by NMC [38]. The size of the depletion observed is smaller but not inconsistent with the DIS results. Remarkably, no enhancement is visible in the $x \approx 0.1$ -0.2 region.

2.4 Summary of the Experimental Situation

We summarize the status of the data on the ratio of bound to free nucleon structure functions. The measurements currently cover the region $5 \times 10^{-5} < x < 0.8$.

- For $5 \times 10^{-5} < x < 10^{-3}$ the only available data are the preliminary ones of the FNAL experiment E665 [41] on the xenon to deuterium ratio. Their statistical accuracy is about $\pm 7\%$; the systematic errors are of similar size.
- In the region $0.0035 < x < 0.35$ the NMC results on He/D, C/D and Ca/D [38] are the most accurate currently available. Statistical and systematic errors are typically 1%. The logarithmic slopes $b = d(F_2^A/F_2^D)/d(\ln Q^2)$ were measured with statistical accuracies $\Delta b \approx \pm 0.02$.
NMC also measured the ratios C/Li, Ca/Li and Ca/Li over the range $0.0085 < x < 0.6$ and $0.8 < Q^2 < 17 \text{ GeV}^2$ [39].
- The best experimental information on the region $0.3 < x < 0.8$ comes from the SLAC E139 data [32] on D, He, Be, C, Al, Ca, Fe, Ag and Au. The statistical errors range from 1 to 3% in each bin, depending on the target nucleus.

- The A dependence in each x bin is approximately reproduced by a function of the type $\alpha + \beta \ln A$.

Precious experimental information is also provided by other types of measurements:

- The NMC results on the ratio of the inelastic and incoherent J/ψ production cross sections from Sn and C targets [56] indicate an enhancement of the gluon distribution in Sn with respect to C. In the region $0.02 < x < 0.2$ the average value of the gluon density ratio is 1.13 ± 0.08 .
- The results from the Drell-Yan experiment E772 [72], sensitive to the anti-quark distribution in the region $0.03 < x < 0.3$, show a shadowing signal similar to that observed in the DIS experiments. It is still an open issue however, from the theoretical point of view, if shadowing in DIS and in Drell-Yan are truly comparable.

The picture that emerges from the data is the following:

1. **Shadowing.** For values of x smaller than about 0.05 the bound to free nucleon structure function ratios are smaller than unity. The size of the depletion appears to depend on the radius as well as on the density of the nucleons. In the region $0.0035 < x < 0.05$ the ratios rise approximately linearly with $\ln x$. No significant Q^2 dependence has been observed in this region. At very small x , for $10^{-5} < x < 10^{-3}$ and for very small values of Q^2 , $0.01 < Q^2 < 0.7 \text{ GeV}^2$, the ratios have been measured to be about constant.
2. **Enhancement.** In the region $0.05-0.1 < x < 0.2$ the structure function ratios become larger than unity. The size of the enhancement is approximately 2% for carbon and calcium. There are indications that the ratio depends weakly on the density, with little or no sensitivity to the nuclear radius. In this region no Q^2 dependence is observed for the helium and carbon to deuterium ratios. The calcium to deuterium ratio shows instead some evidence of a decrease with increasing Q^2 .
In this region the ratio of the gluon distribution in Sn and C as measured in J/ψ production also is larger than unity.
3. **EMC effect.** In the region $0.2 < x < 0.8$ the ratios decrease with increasing x and have a minimum at $x \approx 0.6$. The size of the depletion is approximately proportional to $\ln A$. In this region the ratios show a slight decrease with increasing Q^2 .
4. For $x > 0.8$ the ratios rise above unity again, a consequence of Fermi motion. The experimental information is however very scarce here, the cross section being dominated by quasi-elastic scattering (i.e. elastic scattering on a single nucleon).
5. The integrals of the structure function differences $F_2^A - F_2^D$, evaluated over the range covered by the NMC and by the SLAC E139 data, are compatible with zero. If this conclusion could be extended to the full x range, the implication would be that the momentum carried by quarks is the same in bound and free nucleons.

Chapter 3

Theoretical Review and Comparison with the Data

“All models of the EMC effect are wrong, including mine.” [75]

As we mentioned several times, the discovery of the EMC effect came as a surprise to both the particle and nuclear physics communities: it was generally thought that the only relevant nuclear effects in DIS could be Fermi motion, at large x , and shadowing at very small x and Q^2 . In general the nucleon constituents were believed to be largely insensitive to the nuclear environment surrounding them.

It is fair to mention two exceptions, probably the only *pre*-dictions of the EMC effect. In 1975 Nikolaev and Zakharov [195], addressing the problem of shadowing within a parton model framework, explicitly predicted anti-shadowing for $x \approx 0.05$ -0.1 and stated – in contradiction with a widespread prejudice – that both shadowing and anti-shadowing should not vary with Q^2 . About in the same years, Krzywicki [143], in an attempt to explain some features of inclusive, high p_t pion spectra in proton-tungsten collisions, argued that when the nucleus is probed at very small distances it resembles more closely a single bag of partons than a collection of quasi-free nucleons. Among the consequences of this picture the author mentions “an anomalous nuclear enhancement in the cross section for the deep inelastic lepton scattering at small values of the Bjorken variable”.

It was however only after the presentation of the EMC results on the iron to deuterium structure function ratio that the theoretical community seriously begun to study nuclear effects in DIS. In fact intense theoretical effort has been invested since then and several hundred papers have been produced. Originally the interest focused on the intermediate x region, $0.2 < x < 0.6$, where data were available. It gradually shifted to the region of smaller x , especially after the publication of the EMC NA27 [35] and NA28 [37] results.

In what follows we will attempt to give a survey of the theoretical ideas proposed. We will split the discussion into two parts, examining first the models relevant to the EMC effect in the historical sense of the term, i.e. approximately to the region $0.2 < x < 0.6$, and then those dealing more specifically with shadowing. Although there is some arbitrariness in this distinction, it reflects to a certain extent a difference in the physics on which the models are based.

Several reviews of the theoretical situation exist in the literature. We quote here the ones by Berger and Coester [76], Frankfurt and Strikman [77, 78], Barone and Predazzi [79]. A recent summary of the theoretical situation on shadowing can be found in [78] and [80].

The models of the EMC effect essentially try to explain the apparent softening of the quark distribution in bound nucleons with respect to free ones. They can be broadly classified into two categories:

1. "Conventional nuclear physics" models, describing the effect of the nuclear potential on the nucleons by means of a reduced effective nucleon mass, which implies (since $x = Q^2/(2M\nu)$) a shift of x to higher values and thus a softening of the valence distribution (x *rescaling* or *binding models*). In these models the mass shift is often accompanied by an increased density of virtual pions associated with the nuclear force (*pion models*).
2. Models which require an increase of the quark confinement size in nuclear matter. Because of the uncertainty principle this translates into a reduction of the quark Fermi momentum – and hence of the width of the quark momentum distribution. An increase of the confinement size can be viewed, in the language of QCD, as a rescaling of Q^2 (Q^2 *rescaling*). In some cases a simple increase of the nucleon radius is assumed (*nucleon swelling*). In others quark *deconfinement* is invoked, with the disappearance of the nucleon degrees of freedom in the nucleus in favor of *multi-quark clusters* or even quark-gluon plasma.

In the shadowing region the observed depletion of the nuclear to nucleon structure function ratio is generally explained by two types of models:

1. Generalized Vector Meson Dominance (*GVMD*) models in which the virtual photon is described as fluctuating from a bare photon state to a superposition of vector mesons. These vector mesons interact hadronically with the nucleus and are absorbed mainly by the nucleons on the surface of the nucleus. The inner nucleons are thus "shadowed" by the surface ones and the measured cross section per nucleon is reduced with respect to the free nucleon one.
2. *Partonic models* in which – at least in their original formulation – shadowing is ascribed to the fact that in the Breit frame low x partons, because of the uncertainty principle, spread over a large longitudinal distance. Partons of different nucleons may thus overlap in space and fuse, thereby reducing the density of low momentum partons and increasing that of higher momentum ones.

Figure 3.1 shows the *approximate* range of validity of the various classes of models.

In spite of the theoretical efforts made to understand the physics behind the EMC effect and shadowing there is as yet no unique and generally accepted model of these phenomena. Even worse, the models available are in general of modest predictive power and appear difficult to falsify: most often they just offer a description rather than an explanation of the experimental findings. To the high energy physicist spoiled by the successes of the Standard Model this may appear as an awkward situation; it reflects the difficulty of treating a problem which involves many degrees of freedom and in addition is of non-perturbative nature.

Before turning to a more detailed discussion of the models, one more remark is in order. It is generally assumed that, the deuteron being very weakly bound, F_2^D represents the structure function of an ideal isoscalar nucleon and that $F_2^D = 1/2(F_2^p + F_2^n)$. This assumption, untested for some time after the discovery of the EMC effect, appears now as not so well founded. Effects of a few per cent may be expected both at low x and in the dip region, around $x \approx 0.6$ (see e.g. [81, 82, 136, 137, 168, 227]).

3.1 Models of the EMC Effect

The EMC effect indicates that in the region $0.2 < x < 0.6$ the valence quark distribution in bound nucleons is depleted with respect to that in a free nucleon. Several mechanisms have been proposed to produce such softening.

Models based on “conventional” nuclear physics invoke an excess of virtual pions in the nucleus, associated with the nuclear force. The virtual photon can therefore scatter not only on a quark in a nucleon, but also on a quark (or anti-quark) in a pion. Pions can carry a fraction of the nucleon momentum up to $x = M_\pi/M$. An enhancement of the cross section in the region $x < M_\pi/M$ is thus expected, reflecting the increase in the density of quarks and anti-quarks due to the pion’s constituents. The momentum carried by the pions is lost by the nucleons, and thus by their quarks, which are then on the average slower than if there were no pions. This reduction in the quark momentum may lead to the measured EMC effect pattern.

Attempts have been made to compute the momentum fraction transferred to the pions. This is done by expressing the effect of the nuclear medium in terms of the attractive potential in which the nucleons move: in a nucleus the total energy of a nucleon is reduced with respect to the free particle case because of the (negative) contribution of the potential energy. This can be viewed as a reduction of the effective mass of the nucleon, which in turn implies a shift of $x = Q^2/(2M\nu)$ to higher values (“ x rescaling”). A given virtual photon thus probes the quark distribution in a bound nucleon at larger x than in a free one.

Alternatively, the valence distribution can be made softer by assuming that the quark confinement region increases in size. Because of the uncertainty principle this determines a reduction of the quark Fermi momentum. The width of the valence quark momentum distribution is proportional to the quark Fermi momentum, and is thereby reduced. The valence quark momentum distribution is peaked around $x = 1/3$; its narrowing implies a reduction of the valence quark density at $x > 1/3$, which is what is seen in the data.

In terms of QCD a change in the confinement size means a change of the scale μ^2 appearing in the Altarelli-Parisi equations 1.12-1.13; so the effective value of Q^2 for a bound nucleon is different from – in fact smaller than – that in a free nucleon. This is referred to as “ Q^2 rescaling”. Within this framework the QCD evolution in a bound nucleon “starts earlier”; consequently the amount of sea generated by QCD radiative processes per nucleon is larger in a nucleus than in a free nucleon.

In some models the nucleonic structure tends to disappear at large Q^2 , with quarks and gluons (or a fraction thereof) no longer confined to specific nucleons but rather spreading over the whole nuclear volume, which becomes some sort of a “color conductor”. A similar viewpoint is shared by cluster models, in which quarks are conjectured to move quasi-freely inside enlarged bags, consisting of 6, 9, 12 or more valence quark clusters. It is worth mentioning that both nucleon swelling and multi-quark clusters are not new concepts, but have actually long been advocated by some nuclear physicists.

All these mechanisms give a fair description of the data in the region $0.2 < x < 0.6$, leaving to models of shadowing the task of accounting for the very small x region and ascribing to Fermi motion the rise of the structure function ratio at large x . The very high x region ($x > 1$) has also been the subject of considerable theoretical effort, notably in the framework of partial deconfinement and cluster models, where it is possible for a quark to carry a fractional momentum larger than unity.

3.1.1 Nuclear Binding, Pions and x Rescaling

Shortly after the publication of the EMC results [5], Llewellyn Smith [83] and Ericson and Thomas [84] showed that the gross features of the data could be explained by an enhancement of the pion field in nuclei, associated with the nucleon-nucleon interaction. This point of view was later expanded upon by other authors [85]-[87].

The structure function of a bound nucleon can then be written as the sum of two terms:

$$F_2^A(x, Q^2) = \int_{y \geq x} dy f_N(y) F_2^N(x/y, Q^2) + \int_{y \geq x} dy f_\pi(y) F_2^\pi(x/y, Q^2). \quad (3.1)$$

The first integral is the convolution of the free nucleon structure function F_2^N with the nucleon distribution function in the nucleus f_N . Similarly, the second term is the convolution of the pion structure function F_2^π with the pion distribution function in the nucleus f_π . The assumption is made that the structure functions $F_2^{N,\pi}$ of nucleons and pions are unaffected by the nuclear medium.

In the deep inelastic scattering on a pion at rest in a stationary nucleus x is bounded from above: $x < M_\pi/M \approx 0.15$. We therefore expect that if there are more pions in bound than in free nucleons, their direct contribution (the second integral in equation 3.1) is important at x smaller than ≈ 0.15 . Such a model implicitly leads to an enhanced anti-quark distribution in nuclei, because of the the pion's valence anti-quark.

Furthermore, since these pions carry a fraction of the nucleus momentum, the nucleons are on the average slower than if there were no pions. Let η be the fraction of the nucleus momentum per nucleon carried by pions; then nucleons carry only a fraction $\bar{z} = 1 - \eta$. For a free nucleon this fraction is unity.

The model successfully describes the experimental data in the range $0.1 < x < 0.8$. As an example the number of extra pions per nucleon needed to reproduce the EMC effect in aluminum is about 0.09; they carry approximately 5% of the nucleus momentum. These numbers turn out to be proportional to the nuclear density and rise to 0.11 and 6%, respectively, in gold. Figure 3.2 shows the predictions for the Al, Fe and Au to D structure function ratios compared with the original effect measured by EMC NA2 [5] and with the SLAC E139 [31] and BCDMS [33] results. The model is however less successful in describing the Drell-Yan data [72], since it predicts a strong enhancement of the anti-quark distribution – and thus of the cross section ratio – which is not observed in the data (cf. fig. 2.28; the theoretical curves in the figure are based on the calculations of [88]).

A straightforward extension of the model considers not only pions but also Δ 's [89, 90]: the nucleus is then viewed as a set of nucleons, pions and Δ 's. The structure function F_2 of a Δ baryon is expected to be shifted towards lower x values with respect to the nucleon one [91].

If the Δ 's were the only non-nucleonic component in the nucleus, a fraction of 10-15% such baryons per nucleon would be necessary to reproduce the observed EMC effect in iron for $x > 0.2$ [89]. Such a fraction is somewhat larger than allowed by current nuclear physics estimates [85]. If the nucleus is instead assumed to contain both pions and Δ 's (beyond nucleons), then the effect measured in iron can be reproduced assuming a 4% and 12% admixture of Δ 's and pions, respectively. Figure 3.3 shows the predictions of the model for Al and Fe.

While in the approaches outlined above $\bar{z}$ and the average number of pions in the nucleus are chosen so as to obtain the best agreement with the experimental results, Akulinichev *et*

al. [92, 93] attempted to derive them in a nuclear physics framework. A very similar approach was proposed by Birbrair *et al.* [94]. These authors express the effect of the nuclear medium in terms of the attractive potential V in which nucleons move. The nucleon energy reads in this case:

$$E_N = M + V + \frac{p^2}{2M}, \quad (3.2)$$

where p is the modulus of the three momentum of the nucleon. The effect of the potential V (< 0) can be subsumed into an effective nucleon mass, lower than the free nucleon one, thus causing a shift in the scaling variable $x = Q^2/(2M\nu)$ towards higher values (“ x rescaling”). An earlier discussion of (x dependent) x rescaling can be found in [95]. The same conclusion can be reached from a relativistic extension of the Landau theory of quantum liquids [96]. The bound nucleon structure function F_2^A can be then approximated as:

$$F_2^A(x, Q^2) \approx F_2^D\left(\frac{x}{\bar{z}}, Q^2\right), \quad (3.3)$$

which is the first integral in eq. 3.1 if a δ function for the nucleon distribution $f_N(y)$ is used. The rescaling parameter $\bar{z}$ is a function of V :

$$\bar{z} \approx 1 + V/M \approx 1 + \bar{\epsilon}/M, \quad (3.4)$$

where we have introduced the average one-nucleon separation energy $\bar{\epsilon}$. We recall that the separation energy is the energy necessary to “ionize” a nucleus and make it emit a nucleon. For example, given two nuclei both with Z protons and with N and $N - 1$ neutrons, respectively, their masses are related by the following relationship:

$$M(Z, N) = M(Z, N - 1) + M_{\text{neutron}} - \epsilon_{\text{neutron}}. \quad (3.5)$$

In terms of the binding energies of the two nuclei in question, equation 3.5 can be rewritten as

$$\epsilon_{\text{neutron}} = B(Z, N) - B(Z, N - 1). \quad (3.6)$$

In equation 3.4 $\bar{z}$ is again the fraction of the nucleus momentum carried on the average by the individual nucleon. The nucleon $(Z, N - 1)$ can in principle be in an excited state, so that $B(Z, N - 1)$ need not be the ground state binding energy, thus leading to values of the separation energy considerably larger, in absolute value, than those normally quoted in the literature. For instance for iron $\bar{\epsilon}_{Fe} = -39$ MeV (corresponding to $\bar{z}_{Fe} = 0.96$) is needed in order to reproduce the experimental results (typical values of $\bar{\epsilon}$ are ≈ -10 MeV). Figure 3.4 compares the predictions of this model to the SLAC E139 results [31].

The present approach has been criticized for two reasons.

1. Frankfurt and Strikman [97] pointed out that the treatment of Akulinichev *et al.* uses an incorrect normalization of the relativistic spectral function and leads to an overestimate of the nuclear binding contribution to the EMC effect by up to a factor of three (the so-called flux factor). A similar case was made by Jung and Miller [98]. Analogous

criticisms apply to other relativistic treatments of the nucleus, like the one of Morley and Schmidt [99]. There are approaches at variance with that of Frankfurt and Strikman, e.g. the one proposed by Molinari and Vagrado [100]; these authors make however no explicit prediction for the EMC effect.

2. Li, Liu and Brown [101] showed that the separation energies used by Akulinichev *et al.* are too large (in absolute value) as they do not include an appropriate treatment of density-dependent interactions. In a Hartree-Fock description of the nucleus, they find for iron $\bar{\epsilon}_{Fe} = -26$ MeV and $\bar{z}_{Fe} = 0.972$; this further reduces the role of binding in explaining the EMC effect by an extra factor of two¹.

By means both of the relativistic treatment à la Frankfurt and Strikman and of the separation energies of Li, Liu and Brown, it is found that nuclear binding may account by up to about 20% only of the observed effect.

Kulagin [104] later revisited the calculations of [92, 93] in a relativistic covariant framework, with appropriate treatment of the flux factor. The rescaling parameter should now read

$$\bar{z} = 1 + (\bar{\epsilon} + 2\bar{T})/M, \quad (3.7)$$

where $\bar{T}$ is the mean kinetic energy of the bound nucleons. This calculation is coupled with an evaluation of the mesonic excess in nuclei, necessary – as we discussed – to preserve momentum conservation. A fair description of the data (for $x > 0.1$) can be obtained only with rather large values of the separation energy, $|\epsilon| \approx 50$ MeV. A similar point on the necessity of taking into account the interaction of the virtual photon with the mesons in the nuclei was made by Birbrair *et al.* [105].

Ciofi degli Atti and Liuti [106, 107] showed that nucleon-nucleon correlations induced by realistic interactions indeed strongly increase the average value of the removal energy. Similar conclusions were reached by Antonov *et al.* [108] in the framework of the coherent density fluctuation model [109], in which nucleon-nucleon correlations and binding effects are taken into account.

The discussion on the contribution of nuclear binding to the EMC effect has not yet settled. Dieperink and Miller [110] studied the problem using realistic nuclear forces [111]. They also find that – at least for infinite nuclear matter – there are large deviations from the mean field picture due to nucleon-nucleon correlations, leading to increased nuclear separation energies. Their conclusion however is that nuclear binding can account for only a fraction ($\approx 70\%$) of the effect measured in large nuclei. Nakano [112] investigated how the EMC effect in the x rescaling approach is related to the symmetry properties of the nucleon distribution function and advocated increased removal energies – or non-conventional ingredients in the nucleus – in order to fully explain the EMC effect.

The convolution formalism itself has been the object of several critical studies [113]–[116]. It is fair to say that there is still little consensus on the details of its implementation. Figure 3.5 shows the results of a recent calculation by Gross and Liuti [116], based on a fully covariant derivation of the convolution formula; the most significant difference with respect to previous derivations is the explicit dependence of the bound nucleon structure functions on the momentum of the nucleon itself. The nuclear input is the same as in [106, 107].

¹These points in turn raised some critical discussion, see e.g. [102] and, more recently, [103].

Before closing this section we recall that the nucleon mass shift and the enhancement in the pion density in nuclei have also been given a somewhat unconventional theoretical justification by Preparata and Rattcliffe [117, 118]. According to their theory, the nucleons in the nucleus are involved in a coherent type of dynamical evolution in which all the individuals inside domains of size ≈ 4.5 fm oscillate in phase with a condensed pion field.

3.1.2 Q^2 Rescaling Models

The starting point of the Q^2 rescaling model, formulated by Close, Roberts, Ross and Jaffe [119]-[121] is the observation that

$$F_2^A(x, Q^2) \approx F_2^D(x, \xi Q^2), \quad (3.8)$$

where $\xi > 1$. This suggests that the effective value of Q^2 may not be the same for a nucleus and for a nucleon.

We recall from section 1.1.3 that in QCD Q^2 always appears scaled by either Λ^2 or μ^2 . At $Q^2 \approx \mu^2$ gluon radiation can be neglected and the nucleon appears as made up of the three valence quarks only; the quantity $\lambda \sim 1/\mu$ can then be taken as an estimate of the valence quark confinement size. One can suppose that the quark confinement size in a nucleus, λ_A , is larger than that in a free nucleon, λ_N . This idea was first proposed by Jaffe [122]; its relevance to the EMC effect becomes obvious if the uncertainty principle is considered, as we discussed above. The scales μ_A^2 and μ_N^2 are then also different, and $\mu_A^2 \sim \mu_N^2 \lambda_N / \lambda_A$. The effective value of Q^2 for bound nucleons is thus smaller than that for free nucleons: Q^2 needs to be rescaled to higher values in order to obtain the quark distributions in a free nucleon from those in a bound one.

From the arguments just presented it is clear that the rescaling factor ξ must be a function of μ_N / μ_A ; the actual form can be derived by using the structure function moments M_n defined in formula 1.11. One can start at $Q^2 = \mu^2$, where the valence picture is a good approximation. So for $Q^2 = \mu_N^2$ in a nucleon and for $Q^2 = \mu_A^2$ in a nucleus, the quark distributions are equal and so are the moments: $M_n^N(\mu_N^2) = M_n^A(\mu_A^2)$. If Q^2 changes, the moments of the nucleon structure functions evolve according to the QCD prediction:

$$\frac{M_n^N(Q^2)}{M_n^N(\mu_N^2)} = \left[\frac{\alpha_s(Q^2)}{\alpha_s(\mu_N^2)} \right]^{d_n}, \quad (3.9)$$

where d_n is the QCD anomalous dimension. For bound nucleons a similar relation must hold:

$$\frac{M_n^A(Q^2)}{M_n^A(\mu_A^2)} = \left[\frac{\alpha_s(Q^2)}{\alpha_s(\mu_A^2)} \right]^{d_n}, \quad (3.10)$$

hence, using $M_n^N(\mu_N^2) = M_n^A(\mu_A^2)$,

$$\frac{M_n^A(Q^2)}{M_n^N(Q^2)} = \left[\frac{\alpha_s(\mu_N^2)}{\alpha_s(\mu_A^2)} \right]^{d_n}. \quad (3.11)$$

The rescaling parameter ξ is now defined by the relationship $M_n^A(Q^2) = M_n^N(\xi(Q^2)Q^2)$; solving it to leading order in α_s , one finds that ξ varies with Q^2 but, rather remarkably, is independent of n :

$$\xi(Q^2) = \left(\frac{\mu_N^2}{\mu_A^2} \right)^{\alpha_s(\mu_A^2)/\alpha_s(Q^2)} \quad (3.12)$$

At $Q^2 \approx 20 \text{ GeV}^2$ the Fe/D data indicate a value of the rescaling parameter $\xi \approx 2$. This translates into $\lambda_{Fe} \approx 1.15\lambda_N$: the quark confinement radius for iron thus appears to be 15% larger than for a free nucleon. For helium the prediction is $\lambda_{He} = 1.079\lambda_N$ and $\xi(Q^2 = 20 \text{ GeV}^2) = 1.43$.

So far the model provides no “explanation” for the observed change of scale, only a framework for discussing it. The physics of the change in the confinement radius may lie in nucleon overlap, which makes it more likely for quarks to leak outside the confinement radius of the free nucleon; in this case the effect is controlled by the nuclear density and by the nucleus surface to volume ratio. The nucleon-nucleon overlap probabilities can be estimated by using the nuclear density and the nucleon-nucleon correlation function. Figure 3.6 compares the SLAC E139 results [31] with the model predictions thus obtained. In the range of applicability of the model ($0.2 < x < 0.8$) the agreement with the data is good.

The Q^2 dependence of the EMC effect in the Q^2 rescaling model was discussed by Bickerstaff and Miller [123]. Except at large x , where Fermi motion dominates, very little variation of the structure function ratio is predicted even over a wide Q^2 range. This is a consequence of the fact that the difference between the bound and free nucleon structure functions is small and their evolution slow. The Q^2 dependence predicted by this model is similar to that expected within the framework of nuclear binding and pion models. Figure 3.7 shows the predictions for F_2^{Fe}/F_2^D .

The QCD based Q^2 rescaling model just described and the conventional nuclear physics approach of the x rescaling model discussed before, although rather different in their starting points, can be formally connected [124, 125]. The average nucleon separation energy $\bar{\epsilon}$ and the Fermi momentum k_F can be related to the QCD anomalous dimensions d_n and to the increase of the confinement size. The fact that the nucleon, in the x rescaling model (disregarding pions), carries only a fraction $\bar{z} \approx 1 + \bar{\epsilon}/M$ of the nucleus momentum can be expressed in terms of the structure functions moments as

$$M_{n=2}^A = \bar{z} M_{n=2}^N. \quad (3.13)$$

By comparing this relation with equation 3.11, one obtains:

$$\bar{\epsilon} = M \left[\left(\frac{\alpha_N}{\alpha_A} \right)^{d_2} - 1 \right]. \quad (3.14)$$

Proceeding to the higher order moments one finds for $n = 3$:

$$k_F^2 = 5M^2 \left[\left(\frac{\alpha_N}{\alpha_A} \right)^{d_3} - \left(\frac{\alpha_N}{\alpha_A} \right)^{2d_2} \right], \quad (3.15)$$

where a Fermi gas distribution has been assumed for the nucleons in the nucleus. Equation 3.15 relates the Fermi momentum k_F to the confinement size and the parameters of QCD. Using $n = 4$ yields new relations for $\bar{\epsilon}$ and k_F which are consistent with those given above.

This connection between Q^2 and x rescaling is interesting as it allows to express nuclear physics quantities in the language of QCD. The physical basis for this remains however unclear [126]. Dunne and Thomas [127] in fact argued that x and Q^2 rescaling, instead of being two different facets of the same physical process, should rather be viewed as complementary explanations of the EMC effect². While x rescaling is still seen as a consequence of nuclear binding, Q^2 rescaling is assumed to take into account the fact that the structure function F_2^N in equation 3.1 is that of a nucleon off the mass shell: the dependence of F_2 on the invariant mass of the target nucleon is given by

$$\frac{M_n^A(Q^2)}{M_n^N(Q^2)} = \left[\frac{\alpha_s(P_N^2)}{\alpha_s(P_A^2)} \right]^{d_n}, \quad (3.16)$$

which is the same as equation 3.11 except that μ_A^2 and μ_N^2 have been replaced by the invariant masses of the struck nucleons in the two targets, respectively. The change of scale is thus a change of the mass and not of the confinement scale as in the original Q^2 rescaling model.

An alternative approach to Q^2 rescaling is the one of Fredriksson [128], who described the nucleon as being mostly in a bound quark-diquark state, with the diquark $(ud)_0$ a small, spin zero, tightly bound quark pair. Fredriksson suggested that the diquark may be somewhat bigger in a bound than in a free nucleon as a consequence of the perturbing forces exerted by the quarks of the nearby nucleons. The effect of a swelling diquark on the structure functions is included by scaling Q^2 with the diquark mean square radius. The structure function of a bound nucleon thus reads

$$F_2^A(x, Q^2) \approx F_2^D(x, kQ^2), \quad (3.17)$$

where now $k = \langle r_d^2 \rangle_A / \langle r_d^2 \rangle_D$, with $\langle r_d^2 \rangle$ the diquark radius squared. This is a consequence of the fact that the diquark form factor is expected to scale with $\langle r_d^2 \rangle Q^2$. The Q^2 rescaling parameter is therefore different from the one proposed by Close, Roberts, Ross and Jaffe and, to first approximation, is simply given by the ratio squared of the bound to free nucleon confinement radii. The agreement of the model with the data is acceptable between $x = 0.2$ and $x = 0.6$. For iron $k \approx 1.2-2$, corresponding to an increase of the diquark radius of 10-45%. A critical discussion of this model compared with the one of [119]-[121] is presented in [129].

3.1.3 Nucleon Swelling

That a small change of the nucleon radius, regardless of its origin, can lead to the EMC effect was also shown by Hendry, Lichtenberg and Predazzi [130]. This assumption was later combined with that of a reduced effective mass of bound nucleons with respect to free ones (a concept common to x rescaling models) [131]. The resulting predictions are in reasonable agreement with the SLAC E139 data [31] for $x > 0.2$. It is found that the nucleon radius is about 5 to 10% larger in iron than in deuterium; the effective mass is taken to be 5% smaller

²This approach was however criticized by Frankfurt and Strikman [97] on the grounds of an incorrect normalization of the relativistic spectral function.

than the free nucleon one. Interestingly, in a later paper [132] the authors remarked that also at small x the behavior of the nuclear to free nucleon structure function ratio (shadowing, see next section) can be ascribed to nucleon swelling.

The concept of nucleon swelling is older than the discovery of the EMC effect. A 30% increase in the nucleon size in Fe was for instance advocated [133, 134] to account for quasi-elastic electron scattering data which revealed a softer charge form factor in bound nucleons than in free ones. Already at that stage nucleon swelling was related to a rescaling of the nucleon mass; this connection is automatic in the MIT bag model [135].

Frankfurt and Strikman [77, 136] also argued in favor of nucleon swelling in nuclei in order to describe the EMC effect for $x > 0.3$. In QCD a hadron can be described as a superposition of quark-gluon configurations of different sizes. For example at small x large effective size configurations dominate; in the nuclear physics language this corresponds to saying that the pion cloud of the nucleon extends to large distances. At larger x (> 0.3) and for $Q^2 > 1 \text{ GeV}^2$) configurations with only the three valence quarks dominate, with no pion clouds, sea quarks or gluons. Their size is smaller; the interaction of a bound nucleon in such a configuration with other nucleons is weaker than that of larger size configurations. This property is often called “color transparency”: in the limit of a strictly pointlike hadron, the color field visible outside the hadron is zero. Based on energy considerations, these small size, high momentum configurations are suppressed in bound nucleons with respect to free ones. The suppression is related to the average kinetic energy of a nucleon in a nucleus and to the average binding energy per nucleon. The corresponding increase in the nucleon radius is about 2-4%, smaller than that found in other models.

A very recent formulation of nucleon swelling is the one of [137]. The nucleon is treated as a bound state of three constituent quarks; a simple gaussian ansatz is used for the constituent quark wave functions. Glue and sea quark distributions are derived from the valence one by means of the QCD evolution equations; an approximate procedure is used for the $Q^2 < 0.5 \text{ GeV}^2$ region, where the standard evolution equations are not applicable. The model has one free parameter, $w_q = Rm_q$, where R is the radius of the nucleon and m_q is the constituent quark mass. The value of w_q is fixed to ≈ 0.1 by comparison with the structure function data at large x . The resulting valence, gluon and sea distributions for the nucleon are in good agreement with the data.

The amount of swelling – or stretching, as these authors prefer to name it – that a bound nucleon undergoes is, as usual, taken to be proportional to the nucleon overlapping volume. The increase in the nucleon radius ranges from 4% in He to 9% in Pb. The model takes into account Fermi motion as well as nuclear binding; the latter implies some sort of x rescaling, in that, the nucleus mass M_A being smaller than AM_N (where M_N is the nucleon mass), the appropriate scaling variable is $x_A = AQ^2/(2M_A\nu)$ rather than the usual x . Figure 3.8 compares the predictions of this calculation with the SLAC E139 results [32].

This description of the EMC effect is combined with a treatment of shadowing; we shall therefore come back to it later on, at the end of section 3.2.4.

3.1.4 Deconfinement and Color Conductivity

Q^2 rescaling is a natural by-product of the color conductivity model [138]-[140]. The basic idea of this approach is that when a nucleus is observed with high resolution, i.e. at large Q^2 , the nucleonic structure becomes unimportant: quarks and gluons are no longer confined to specific

nucleons but rather spread over the whole nuclear volume. At large Q^2 the nucleus thus turns into a medium in which colored objects can travel freely, a “color conductor”. In practice the effective nucleon radius, i.e. the quark and gluon confinement size, is taken to increase with Q^2 until full color conductivity sets in; the basic point of the Q^2 rescaling model [119]–[121], namely the variation of the confinement radius with Q^2 , is thus recovered. With the explicit further assumption that the running coupling constant α_s is independent of A , one indeed finds again eq. 3.12 for the rescaling parameter ξ . The possibility that the QCD scale parameter Λ (and thus α_s) is a function of A is also explored. Instead of eq. 3.12 in this case one finds simply $\xi = (R_A/R_D)^2$ (with R_A , R_D the radius of nucleus A and of deuterium, respectively), a result similar to that of the diquark model by Fredriksson [128]; the corresponding variation of the effective value of Λ is $\Lambda(A_1) = \Lambda(A_2)(R_{A_2}/R_{A_1})$. Figure 3.9 shows the predictions of the model together with the original EMC NA2 data on Fe/D [5] and the N/D data of BCDMS [33] (labelled “This experiment”).

A slightly different viewpoint on deconfinement was expressed by Gupta *et al.* [142]. In their model the influence of the nuclear medium on the parton distributions consists in changing the phase space volume open to quarks and gluons. In the infinite momentum frame the quarks in the nucleus are assumed to have a minimum transverse momentum $p_t^{\min}$ proportional to the inverse of the nuclear radius:

$$p_t^{\min} = \beta A^{-1/3},$$

where $\beta \approx 100\text{--}200$ MeV. Thus the cutoff in p_t decreases with increasing A and the phase space available to quarks increases. The scale parameter μ in the Altarelli-Parisi equations 1.12–1.13 is set in this model equal to $p_t^{\min}$. This in turn means that the *effective* value of Λ_{QCD} extracted from data on nuclear targets, λ_A , is also a function of A :

$$\lambda_A = \Lambda A^{-1/3}.$$

A comparison of this model with the data is shown in fig. 3.10.

Deconfinement has been also approached in a thermodynamical framework, in which a fraction of deconfined partons is treated as a non-interacting Fermi gas at temperature T .

Krzywicki and Furmański [144], developing the ideas presented in [143] (one of the only two partial *pre*-dictions of the EMC effect, together with that of Nikolaev and Zakharov [195]), regarded the whole nucleus as a dilute gas of quarks and gluons. The model is of course somewhat extreme, but reproduces the trend of the data, at least qualitatively, if $T = 30$ MeV. If only a few nucleons are instead assumed to “deconfine”, with their constituents released into the full nuclear volume [145, 146], the agreement with the data improves (cf. fig. 3.11). The temperature of the deconfined parton is again found to be in the range 25–45 MeV and the average number of deconfined nucleons per nucleon is about 0.10–0.17. Figure 3.11 also shows the expected ratio for the gluon distributions in Fe and D. The trend disagrees with the recent NMC data on Sn/C [56].

The interpretation of structure functions in thermodynamical terms is not new. Angelini and Pazzi [147], before the discovery of the EMC effect, treated the valence quarks as a non-interacting, relativistic Boltzmann gas. They found that the measured structure functions could be well reproduced for $x > 0.1$ if the temperature was assumed to be about 50 MeV. In

this framework the EMC effect can be understood in terms of a different quark gas temperature in bound and free nucleons [148]. For iron and deuterium the temperature difference is found to be $\Delta T = 3 \pm 1$ MeV. Furthermore the effective nucleon volume, i.e. the volume occupied by the quark gas, turns out to be 30% larger in Fe than in D, somewhat higher than many of the estimates for nucleon swelling given above.

3.1.5 Cluster Models

A variation on the themes of deconfinement and nucleon swelling is that of clusters. Several authors conjectured that a fraction of the valence quarks in the nucleus move quasi-freely inside enlarged bags, consisting of 6, 9 or more valence quark clusters. The possibility that clusters might explain the EMC effect was first mentioned by Jaffe [122] in the MIT bag framework.

Pirner and collaborators [149]-[151] proposed one such model before the discovery of the EMC effect in order to explain deep inelastic data from ^3He in the region $1 < Q^2 < 4 \text{ GeV}^2$ [152]. In their approach, two nucleons are assumed to start forming a six-quark cluster when they are a distance $2R$ apart, where R is an effective nucleon radius; similarly, if a third nucleon is less than $2R$ from either of these two nucleons then a 9-quark cluster forms. The probability of having a cluster of a given size can then be calculated by assuming realistic nucleon-nucleon potentials or by resorting to Monte Carlo simulations [153].

The momentum distribution of quarks in an i -quark cluster, $n_i(x)$ ($i = 6, 9, \dots$), is taken to be

$$x n_i(x) \sim x^{1/2} [1 - x/(i/3)]^{2(i-1)-1},$$

based on counting rules [154, 155].

The experimental results on ^3He [152] are well described with an effective nucleon radius of 0.5 fm, leading to a 10-20% probability of non-nucleonic quark structures in the nucleus. The success of the quark-cluster model in explaining the ^3He data depends on the fact that a quark in a 6 or 9-quark bag has the possibility of carrying the momentum of two or three nucleons, thus providing high momentum components not present in a conventional picture based on 3-quark systems only.

Using similar arguments, Carlson and Havens [156] showed that the EMC effect in Fe can be described by assuming that nucleons are, with a 30% probability, subsumed into 6-quark or larger clusters. A similar discussion of 6-quark clusters has been proposed by Chemtob and Peshanski [157].

Faissner *et al.* [158]-[160] concentrated on the role of 12-quark clusters, i.e. α particles, which are known to be very stable systems. Their stability is a consequence of the large binding energy and can be easily understood in the QCD framework. We can picture a nucleon as made up of three quarks connected by color strings. Let us assume that the three quarks sit at the vertices of an equilateral triangle. While it is difficult to form a stable system out of two or three of these triangles, it is easy if four are considered, as they form a regular tetrahedron. The quarks thus loose their assignment to individual nucleons: the three differently colored quarks at each corner of the tetrahedron may be combined into a white nucleon just as well as those belonging to the original triangles. The α particle is therefore not just a 4-nucleon compound, but also a system of 12 valence quarks which form a singlet in spin, isospin and color.

In order to evaluate the structure functions of an α -cluster, a model is adopted in which the incoming lepton scatters off a pointlike constituent quark with an effective mass $m_q(Q^2)$. The motion of the constituent quark inside the cluster is described in terms of a distribution function f_q and the deep inelastic scattering cross section is obtained as the convolution of the lepton-quark cross section folded with the quark distribution f_q ; each quark is taken to give an incoherent contribution.

The EMC effect data can be described by assuming that the Fermi momentum of the quarks inside the clusters is lower than inside the nucleon and that there is a fraction g_{cl} of α clusters in the nucleus. A lower Fermi momentum simply means that the size of the cluster is larger than the nucleon one. The EMC data can be reproduced in the region $0.2 < x < 0.6$ with a clustering probability $g_{cl} \approx 10\%$, with an effective quark mass $m_q = 150$ MeV and with a cluster radius equal to that of the α particle. The Fermi momentum of the quarks inside a free nucleon is set to 300 MeV and the explicit form of the distribution function f_q is found by assuming a harmonic oscillator potential for the quarks inside the cluster.

In this framework the A dependence of the EMC effect has a geometrical origin: clusters are more likely to form well inside the nucleus rather than on its surface, where the number of neighbors is smaller. Figure 3.12 shows the results thus obtained. Similar arguments on the geometrical origin of the A dependence of the EMC effect in a multi-quark cluster model were proposed by Daté *et al.* [161].

The relevance of α clusters was also discussed by Kondratyuk and Smatkov [162] who found that a 15-20% admixture of 12-quark clusters in the nucleus may explain not only the EMC effect data but also the behavior of the ^4He form factor at large Q^2 . The momentum distribution of the quarks in the cluster is taken to be exponentially falling, $\sim x \exp(-ax)$, with $a \approx 8-10$; this form is based on the assumption that the quarks inside a cluster behave like a Fermi gas.

A similar fraction of α clusters is found in another model calculation by Clark *et al.* [163]. The authors remark however that the experimental results can be reproduced just as well with any two-component cluster model.

A dynamical mechanism that can lead to clusters of three and four nucleons was proposed by Barshay [164, 165] and is depicted in fig. 3.13. When three nucleons are very close to each other, a valence quark from each nucleon may be exchanged to each neighbor with the emission of a hard gluon. The three gluons interact and the resulting repulsive force tends to give to the system an equilateral triangle configuration. The emission of these gluons softens the momentum distribution of the valence quarks and this leads to the EMC effect.

Quantitatively this idea is implemented by assuming that a fraction δ of the valence quarks moves in a larger bag with a gaussian distribution of the type $Cx^{1/2} \exp(-Bx^2)$, where the normalization C is chosen in such a way as to keep the number of valence quarks constant and the parameter B reflects the spatial extent of this motion, about 1.3 fm. The structure function of a bound nucleon thus reads:

$$F_2^A(x) = \frac{5}{18} [xV(x)(1 - \delta(A)) + Cx^{1/2} e^{-Bx^2} + xS(x)], \quad (3.18)$$

where $V(x)$ and $S(x)$ are the valence and sea quark distributions in a free nucleon. The A dependence of δ is taken to be the same as determined in [161], with the nucleons on the surface of the nucleus excluded from the three-nucleon correlation.

The mechanism can be thought to work also for sea quarks and anti-quarks, thus leading to a description of the shadowing region as well. In this case the distribution of the sea quarks

involved in the three-body correlation is assumed to be of the type $C'x^\beta \exp(-B'x^2)$. The parameters B , B' , β and δ can be estimated only approximately. What is done then is to fix them with a fit to the data, which renders the predictive power of the model rather modest, especially at small x .

With increasing Q^2 , the valence quarks involved in the correlation will undergo an additional softening as a consequence of QCD evolution. This leads, for $x > 0.25$ to a further decrease of the structure function ratio. Figure 3.14 shows the results for Ca. At $x = 0.15$ the predicted Q^2 dependence is qualitatively in disagreement with the NMC results (cf. fig. 2.15).

The model makes explicit predictions for the ratio of the glue distribution in bound and free nucleons as a function of ν and x for different z and p_t^2 ranges. The ratio is expected to have a maximum of a few per cent above unity, 2 to 5% for the Ca/D case depending on the input parameters, for x between 0.1 and 0.4, again depending on the input parameters. The prediction is in qualitative agreement with the NMC results [56], which are however for the Sn/C ratio.

Hoodboy and Jaffe [166] offered another view of the formation of quark clusters, which they also describe in terms of quark exchange between nucleons. The calculation is carried out for $A = 3$ nuclei only, for which reliable wave functions are available; a nuclear wave function completely anti-symmetrized in the nuclear coordinates as well as in the quark coordinates is used. Quark exchanges between two nucleons at the time are considered only. The predictions of the model are similar to those of the Q^2 rescaling model [119] in the range $0.2 < x < 0.6$.

Lassila and Sukhatme [167] proposed a quark-cluster model in which the contribution of 6-quark clusters only is included. The valence and sea quark distributions in a cluster with N valence quarks are taken to have the forms:

$$\begin{aligned} V_N(x) &\sim x^{1/2}(1-x)^{b_N} \\ S_N(x) &\sim (1-x)^{a_N}. \end{aligned}$$

The exponent b_N is set to $b_N = 2N - 3$ (corresponding again to the simplest application of the dimensional counting rules [154, 155]); this gives $b_6 = 9$. The value found for $N = 2$, agrees with the measured pion structure functions. The sea distribution exponent is taken for a 6-quark cluster to be $a_6 = 11$; the results do not however depend strongly on this choice. Figure 3.15 shows the structure function ratio thus found (curve "A"); if one takes $b_6 = 10$ the curve labelled "B" is obtained. Finally curve "C" corresponds to $a_6 = 13$. The model predictions do not show any enhancement for $x \approx 0.1-0.2$; the authors claim however that the inclusion of 9-quark clusters causes the ratio to become larger than unity in this region. The predictions are compared with the EMC NA2' data [35] on carbon, copper and tin, and with the EMC NA2 [5], BCDMS [33] and SLAC E87 [27] results for iron.

The same authors also made explicit predictions for the ratios of the gluon distributions in bound and free nucleons [168]. Figure 3.16 shows the results, together with the ratios expected for valence and sea. Figure 3.17 compares the ratio of the gluon densities in Sn and C as measured by NMC [56] with the model predictions.

It is worth remarking that this model overlaps to some extent with the parton recombination approach used to describe the shadowing region and discussed below, with the nucleon clustering probability related to the amount of parton recombination. Actually, like parton recombination models, it can reproduce [170] the shadowing signal seen by the Drell-Yan experiment E772 [72] (cf. fig. 3.18); this fact is at variance with the conclusions contained in [72],

based on the calculations of [88]. The predictions vary slowly as a function of x_1 , the fraction of momentum the incident proton carried by the annihilating parton.

Dias de Deus *et al.* [171] combined quark-clusters and Q^2 rescaling, introducing the scale factor $R_i^2 \sim 1/\mu_i^2$ which measures the mean free path of the quarks in the i -cluster; here μ_i^2 is the QCD scale parameter appropriate for the cluster: at $Q^2 \approx \mu_i^2$ the cluster can be viewed as made up by i valence quarks only. For larger values of Q^2 sea quarks, anti-quarks and gluons are generated. Equation 3.12 can now be written in the following form:

$$\frac{Q_i^2}{Q_j^2} = \left(\frac{\mu_i^2}{\mu_j^2} \right)^{\alpha_s(\mu_j^2)/\alpha_s(Q_j^2)}, \quad (3.19)$$

where Q_i^2 and Q_j^2 are the Q^2 values for which the same “amount” of QCD evolution has occurred in the i and j clusters. The relationship between R_i and A is determined according to the ansatz

$$p = \frac{1 + \exp(-a)}{1 + \exp(-aA^{1/3})} \quad (3.20)$$

(where a is a constant and p the average number of nucleons per cluster), while for the dependence of R_i on p , the simple relation $R_i = R_1 p^{1/3}$ is used, with R_1 the confinement scale in deuterium. A fit to the EMC NA2 data [5] yields $p = 1.05$ for iron. In a later paper [172] the same authors repeated the fit using the SLAC E139 results [31]. The ansatz 3.20 is replaced here by the simpler one $p = A^\lambda$, where λ is a free parameter. The fit gives $\lambda = 0.012$, which reproduces the value for $p(Fe)$ previously found. Figure 3.19 shows the results.

In closing this discussion of the cluster models we recall that they offer (together with all models predicting partial deconfinement) a natural framework for studying nuclear structure functions for x values larger than unity, the so-called cumulative region (for a discussion see for instance [10], [151] and [173]-[175]): the cross section ratios in the $1 < x < 2$ region are affected by the existence of 6-quark clusters, in which a quark can carry up to twice the momentum of an individual nucleon. Similarly 9-quark clusters dominate the region $2 < x < 3$. Unfortunately no reliable experimental data exist for $x > 1$, essentially because of the difficulty of extracting the deep inelastic signal from the quasi-elastic one, which dominates in this region.

3.2 Models of Shadowing

The total hadron-nucleus cross section σ_{hA} is known to be smaller than A times the total hadron-nucleon cross section σ_{hN} :

$$\sigma_{hA} < A\sigma_{hN}.$$

Loosely speaking this is ascribed to the fact that the nucleons on the surface of the nucleus project their “shadow”, in an optical analogy, onto the inner ones, which are thus not exposed to the full hadron beam intensity. If the hadron-nucleus cross section is large, then for large A the effective number of nucleons participating in the inelastic interactions is reduced to the number of surface nucleons, $\approx A^{2/3}$.

A theoretical framework describing the interaction of hadrons with nuclei is the Glauber multiple scattering model [176]. The basic assumption of the Glauber treatment is that the amplitude for the interaction of a high energy hadron with a nucleus can be built up from the scattering amplitudes off the individual nucleons. The phase shift experienced by the projectile is simply taken to be the sum of the phase shifts that it would experience by interacting with each of the nucleons in isolation. The motion of the nucleons is neglected and their positions are taken as fixed. Furthermore the assumption is made that the target nucleons do not overlap significantly and that the incident hadron scatters mainly through small angles. In this framework the process in which the incoming hadron interacts with a nucleon and successively reinteracts with a second one turns out to have an amplitude opposite in phase (when the incoming hadron energy is large and the hadronic scattering amplitude is mainly imaginary) to that of the single scattering process, in which the hadron interacts with a single nucleon only. The cross section *per nucleon* is thus smaller in the rescattering case, giving rise to shadowing. For a nucleus with A nucleons also triple, quadruple, ..., A -fold interactions must be taken into account. The successive terms in this multiple scattering series decrease however in size; they also alternate in sign, so that while the double scattering amplitude is negative, the triple scattering one is positive etc.; this feature leads to the familiar diffraction-like pattern of hadron-nucleus cross sections.

High energy, photon-induced reactions bear a remarkable resemblance to purely hadronic reactions, apart from a scale factor $\simeq \alpha$ in the cross section, where α is the electromagnetic coupling constant. This finds a simple interpretation in terms of the hadron structure of the photon [177]-[179]. The physical photon can be viewed (fig. 3.20) as fluctuating between a bare photon state, a set of electromagnetic states (e.g. e^+e^- pairs) and a superposition of hadronic states with the same quantum numbers as the photon ($J^{PC} = 1^{--}$), namely vector mesons; the probability for the photon to convert to a hadron is about α . This is the basis of the Vector Meson Dominance (VMD) model. It is clear then that the photon-nucleus cross section must exhibit some of the characteristics of the hadron-nucleus interaction. Photon interactions should thus also be shadowed in nuclei, as first predicted by Stodolsky [180]. This is indeed observed [179]; the size of the effect increases with the photon energy and with the nucleon mass number A .

The VMD arguments sketched for the real photon case apply to virtual photons as well. The experimental evidence for shadowing of virtual photons was for a long time neither overwhelming nor consistent (fig. 2.2). Some effect was seen at low Q^2 , but considerably less than that found in the photo-production experiments. Shadowing appeared to decrease rapidly as Q^2 increased. The situation became clearer with the EMC NA2' [35, 36] and the EMC NA28 results [37] in which shadowing was unmistakably present in the small x region. In contrast to the earlier data, the effect was there also at values of Q^2 significantly larger than 1 GeV^2 .

Originally only the lowest mass vector states (ρ , ω , ϕ) were considered in VMD models; it is clear now that heavier ones are also necessary in order to reproduce the experimental results. This approach, known as Generalized Vector Meson Dominance (GVMD), is able to give a reasonable description of the data up to $x \approx 0.1$. At least three different formulations of GVMD exist (Schildknecht *et al.* [183]-[188], Piller *et al.* [189]-[191] and Shaw *et al.* [192]-[194]), differing in the details of the implementation of the ideas just outlined, but all reproducing the EMC data.

While the framework in which the shadowing of photons was initially predicted and interpreted was the VMD one, a partonic approach, pioneered by Nikolaev and Zakharov [195], is also available. Shadowing is in this case ascribed to a depletion of partons in the small x

region. In the Breit frame, low x partons spread over a large longitudinal distance because of the uncertainty principle; small momentum partons belonging to different nucleons may thus occupy the same region of space, they may interact and fuse. This leads to a decrease of the number density of low x partons with respect to the free nucleon case and – because of momentum conservation – to an increase in the density of higher momentum ones (“anti-shadowing”). Shadowing and anti-shadowing are not expected to vary appreciably with Q^2 . The predictions of anti-shadowing and of the Q^2 scaling of shadowing and anti-shadowing are peculiar of parton recombination approaches.

The discussion of Nikolaev and Zakharov predates QCD. The mechanism they proposed has been investigated more recently in a perturbative QCD framework. Perturbative QCD does not predict the shape of parton distributions, even less so nuclear ones, but rather their Q^2 evolution. This evolution has been studied in detail by Müller and Qiu [196] and Qiu [197]; their analysis confirms the pre-QCD results of Nikolaev and Zakharov, finding only a weak Q^2 dependence for shadowing.

Some semi-quantitative attempts to describe the x shape of shadowing have also been made in this framework, with special emphasis on trying to predict the A dependence of the point where shadowing sets on (Berger and Qiu [198], Close and Roberts [199]). Models of this kind, based on phenomenological parametrizations of the x dependence of shadowing supplemented by QCD arguments, are now able to give an accurate description of the available data (Zhu *et al.* [202]–[207]).

Partonic models have also been formulated differently. In the laboratory system, where the target nucleus is at rest, the interaction of the virtual photon with the target is described in the language of hadron-hadron diffractive scattering by means of pomeron exchange (Castorina and Donnachie [212, 213], Preparata and Ratcliffe [217] and, to a certain extent, Kwieciński and Badelek [218]). The photon interacts then with a pomeron emitted by a nucleon rather than with the nucleon itself. If the pomeron is taken as a superposition of gluons, the photon may interact with the gluons via photon-gluon fusion processes. Stated differently, the photon may be seen to dissociate into a $q\bar{q}$ pair, with the pomeron interacting with the quarks in the pair (Brodsky and Lu [220]). It is clear that if the $q\bar{q}$ pair is looked upon as a meson, we may recover the vector meson dominance picture discussed above. GVMD and partonic approaches have actually blended in a unitary picture in the recent papers of Frankfurt and Strikman [222, 224], Nikolaev and Zakharov [226, 227] and Barone *et al.* [137]. These authors classify the $q\bar{q}$ pairs according to their transverse size: only the pairs of large transverse dimensions interact with hadronic cross section and give rise to shadowing. Nuclear matter is instead more transparent to small size pairs and the interaction cross section of such pairs with the nucleus is just proportional to the number of nucleons in the nucleus.

We turn now to a description of the available models, classifying them according to their approach being VMD-like or partonic. As we just mentioned, the distinction is somewhat artificial and the two types of treatment may have much in common.

Basically all the models we shall discuss offer a good description of the EMC NA2' [35, 36] and NA28 results [37]. Only a few however are able to reproduce satisfactorily the more recent, higher precision ones from NMC [38, 39]. In the following we shall mostly confine ourselves to a comparison of the model predictions with the latter results.

3.2.1 (Generalized) Vector Meson Dominance Models

Let $q_\mu = (\nu, \vec{k})$ be the four-momentum of the photon and M_V be the mass of the hadronic fluctuation; then the energy of the fluctuation will be:

$$E_V = \sqrt{M_V^2 + k^2},$$

to be compared with the original energy of the photon:

$$\nu = \sqrt{-Q^2 + k^2},$$

where, as usual, $Q^2 = -q^2$ and $k^2 = \nu^2 + Q^2$. The energy difference $\Delta E = E_V - \nu$ between the two states, for large values of ν , is

$$\Delta E = E_V - \nu \approx \frac{Q^2 + M_V^2}{2\nu}. \quad (3.21)$$

The hadronic fluctuation will then extend over a distance (the "coherence length")

$$d(M_V^2, Q^2) = \Delta t \approx \frac{1}{\Delta E} \approx \frac{2\nu}{Q^2 + M_V^2} = \frac{1}{M_V^2} \frac{1}{1 + Q^2/M_V^2} \quad (3.22)$$

and will have a mean free path $l(M_V^2) = 1/[\sigma_V(M_V^2)n_0]$, where n_0 is the nucleon density and $\sigma_V(M_V^2)$ is the cross section for interactions of the hadronic system of mass M_V with the nucleon. Necessary conditions for shadowing are that $d(M_V^2, Q^2) > l(M_V^2)$ and $R_A > l(M_V^2)$, with R_A the nuclear radius. Shadowing due to a fluctuation of mass M_V will thus disappear, at fixed ν , with increasing Q^2 because of the decrease of $d(M_V^2, Q^2)$. Conversely, at fixed Q^2 , shadowing will decrease with increasing x . For $Q^2 = 0$ we recover the real photon case and we see that the length over which the hadronic fluctuation travels is directly proportional to ν . Shadowing of real photons will thus be a growing function of the photon energy.

The assertion that the hadronic fluctuations of the photon are solely constituted by the vector mesons ρ^0 , ω and ϕ defines the VMD model in its simplest form [180]-[182]. The less restrictive assumption that higher mass resonances are also involved leads to the Generalized Vector Meson Dominance (GVMD) model.

In both VMD and GVMD the transverse virtual photon-nucleon cross section reads:

$$\sigma_{\gamma N} = \sum_i \frac{M_{V_i}^4}{(Q^2 + M_{V_i}^2)^2} \left(\frac{e}{f_i} \right)^2 \sigma_{V_i} \quad (3.23)$$

where the sum is over the vector mesons $V_1, V_2, \dots, V_N$, the ratio (e/f_i) is the vector meson-photon coupling and σ_{V_i} is the total V_i - N cross section. The quantity

$$\frac{M_{V_i}^4}{(Q^2 + M_{V_i}^2)^2} = \frac{1}{(1 + Q^2/M_{V_i}^2)^2}$$

is the vector meson propagator squared. The only difference between VMD and GVMD is in the number of vector meson states V_i over which the sum is extended. In order to obtain the vector meson-nucleus cross section, $\sigma_{\gamma A}$, the Glauber multiple scattering series is used.

In the cross section 3.23 the possibility of interference between different vector states belonging (in the VMD sense) to the photon is excluded: such (G)VMD models are known as *diagonal* ones. *Off-diagonal* models include these interference terms; the cross section in this case is:

$$\sigma_{\gamma N} = \sum_i \sum_j \frac{M_{V_i}^2}{Q^2 + M_{V_i}^2} \frac{e}{f_i} \frac{M_{V_j}^2}{Q^2 + M_{V_j}^2} \frac{e}{f_j} \sigma_{V_i V_j}, \quad (3.24)$$

where now $\sigma_{V_i V_j}$ is the total cross section for the process $V_i N \rightarrow V_j N$.

It is clear from relations 3.22-3.24 that the inclusion of higher mass resonances makes shadowing less sensitive to variations of Q^2 ; the inclusion of the off-diagonal transitions also may affect the Q^2 dependence of shadowing. The dependence of σ_V or $\sigma_{V_i V_j}$ on the mass of the vector mesons gives a further degree of freedom. Finally, the sums in 3.23 and 3.24 can be extended to an infinite series of vector states; in this case the equal spacing ansatz $M_{V_i}^2 = M_0^2(1 + i\lambda)$ is often taken for the spectrum of the vector states. Alternatively the sums can be turned into integrals. Formula 3.23 then becomes

$$\sigma_{\gamma N} = \int_{M_{thr}^2}^{\infty} dM_V^2 D(M_V^2) \frac{M_V^4}{(Q^2 + M_V^2)^2} \left(\frac{e}{f(M_V^2)} \right)^2 \sigma(M_V^2), \quad (3.25)$$

where the lower limit of integration is generally set to twice the pion mass squared: $M_{thr} = 2m_\pi$ and where we have introduced the density of vector states per unit mass squared interval: $D(M_V^2) = dN/dM_V^2$. The function D and the coupling (e/f) may be related to the cross section for e^+e^- annihilations into hadrons [177]:

$$\sigma_{e^+e^-} = 16\pi^3\alpha^2 \left[\frac{D(M_V^2)}{f^2(M_V^2)} \right]. \quad (3.26)$$

Using this relation and the ratio

$$R_{e^+e^-}(M_V^2) = \frac{\sigma(e^+e^- \rightarrow \text{hadrons})}{\sigma(e^+e^- \rightarrow \mu^+\mu^-)} \quad (3.27)$$

with $\sigma(e^+e^- \rightarrow \mu^+\mu^-) = 4\pi\alpha^2/(3M_V^2)$, one can rewrite the integral 3.25 as

$$\sigma_{\gamma N} = \frac{\alpha}{3\pi} \int_{M_{thr}^2}^{\infty} dM_V^2 R_{e^+e^-}(M_V^2) \frac{M_V^2}{(Q^2 + M_V^2)^2} \sigma(M_V^2). \quad (3.28)$$

Since $R_{e^+e^-} \sim \text{const.}$, in order for the integral in 3.28 to converge and for $\sigma_{\gamma N}$ to have the usual "scaling" behavior $\sigma_{\gamma N} \sim 1/Q^2$, the cross section $\sigma(M_V^2)$ must vary as $1/M_V^2$. This is somewhat non intuitive; off-diagonal models allow the more natural assumption $\sigma(M_V^2) \sim \text{const.}$

Only GVMD models are flexible enough to reproduce the experimental data both for $\sigma_{\gamma N}$ and $\sigma_{\gamma A}$ (notably their Q^2 dependence). We will now review the available formulations of GVMD limiting our discussion to those proposed or revived after the experimental situation was clarified by the quoted EMC results [35]-[37].

One of the longest lived models is the effectively diagonal one of Schildknecht, originally proposed in 1973 [183]-[188]. In this model the sum over the ρ , ω and ϕ contributions is complemented by the integration over a higher mass continuum; the size of the continuum contribution is normalized at $Q^2 = 0$ to be 22% of the cross section, so as to guarantee the correct normalization of the total photo-production cross section. The cross section $\sigma(M_V^2)$ is assumed to vary as $\sigma(M_V) \sim 1/M_V^2$. This behavior is an effective one and in [185, 186] it is shown to be a consequence of the interference between different vector mesons in a full off-diagonal model in which $\sigma(M_V) \sim \text{const.}$ – as expected from the fact that strong interaction cross sections do not depend much on the mass of the particle.

In figure 3.21 we compare the NMC C and Ca to D ratios [38] with the model predictions (continuous curves). The curve superimposed to the carbon data is for $A = 10$. The carbon prediction is for $Q^2 = 14 \text{ GeV}^2$, while the calcium one is for $Q^2 = 2.8 \text{ GeV}^2$; they are taken from [188].

The NMC C/D data are at the same Q^2 as the model prediction only for $x = 0.175$, at the upper edge of the enhancement region. Since the model predicts an increase of the ratio with increasing Q^2 , the expected values for the C/D ratio in fig. 3.21 should be shifted downward for all $x < 0.175$. Unfortunately in [188] quantitative predictions for the Q^2 dependence are given only for $x = 0.0125$. In this bin the expected Q^2 slope $b = d(F_2^C/F_2^D)/d(\ln Q^2)$ varies from $b \simeq 0.02$ at small Q^2 ($0.75 < Q^2 < 1.36 \text{ GeV}^2$) to $b \simeq 0.04$ at larger Q^2 ($4.6 < Q^2 < 7.5 \text{ GeV}^2$), marginally consistent with the measured value, $b = -0.02 \pm 0.02$. If the predicted Q^2 slope is anyway used, one finds $F_2^C/F_2^D(x = 0.0125) \simeq 0.86$, much smaller than the experimentally measured value of $0.915 \pm 0.009 \text{ (stat.)} \pm 0.007 \text{ (syst.)}$.

The Ca/D data have about the same Q^2 as the theoretical curve ($Q^2 = 2.8 \text{ GeV}^2$) at $x = 0.0175$. Here as well, as in the C/D case, the prediction is lower than the experimentally measured value, as can be seen in fig. 3.21. The expected Q^2 slope at $x = 0.0125$ is $b \simeq 0.05$, in disagreement with the measured value of $b \simeq -0.01 \pm 0.02$. If we nevertheless use the predicted slope to obtain the expected value for the ratio in this bin, we find $F_2^{Ca}/F_2^D(x = 0.0125) \simeq 0.83$, lower than the measured value of $0.859 \pm 0.007 \text{ (stat.)} \pm 0.008 \text{ (syst.)}$.

In order to make the comparison quantitative over a larger x region we deduced Q^2 slopes for this model by comparing the predictions for the EMC NA28 data [37] published in [187] and those – at a different Q^2 – given in [188]. The resulting ratios are plotted in fig. 3.21 as dashed lines. The model does not reproduce the features of the data and systematically underestimates the measurements, except at very small x for Ca.

A similar formulation is the one of Piller, Takayanagi and Weise [189]-[191]. The calculation is based on formula 3.28 where the effective vector meson-nucleon cross section $\sigma(M_V^2)$ is taken, here as well, as $\sim 1/M_V^2$.

In the framework of this model (and the previous one) the shadowing effect is dominated by the low mass vector mesons. This is because the propagation length of these states is large compared to the nuclear size; in addition their mean free path in nuclear matter is smaller than that of the high mass components because of the $1/M_V^2$ behavior of the vector meson-nucleon cross section. The authors point out however that without the inclusion of the higher mass mesons it would not be possible to reproduce the Q^2 dependence of $\sigma_{\gamma N}$.

Figure 3.22 compares the model predictions and the NMC data. The agreement is reasonable, although the model tends to slightly underestimate the ratios.

Figure 3.23 shows the Q^2 dependence expected in this model for the carbon and calcium to deuterium ratios. For Q^2 larger than 1 GeV^2 shadowing decreases nearly logarithmically with increasing Q^2 . At very small Q^2 ($< 1 \text{ GeV}^2$) and $x > 0.03$ the opposite trend is pre-

dicted, reflecting the Q^2 dependence of the coherence length $d(M_V^2, Q^2)$. Shadowing is expected to die out with increasing Q^2 , as in all VMD models. For C/D (Ca/D) the slope $b = d(F_2^C/F_2^D)/d(\ln Q^2)$ is predicted to be $b \simeq 0.05$ (0.07) at $x = 0.01$, $b \simeq 0.03$ (0.03) at $x = 0.04$ and $b \simeq 0.02$ (0.02) at $x = 0.06$. The predictions are consistent with the data for carbon. For calcium they are higher than the measured values.

In the off-diagonal GVMD model of Ditsas and Shaw [193, 194], the virtual photon-nucleon cross section is expressed by means of a relation of the type 3.24, in which the sum is extended to a series of mesons equally spaced in mass; only diagonal and next-to-diagonal transitions are allowed, i.e. $i = j \pm 1$. Furthermore the diagonal vector meson-nucleon cross sections $\sigma_{V_i V_i} = \sigma_{V_i}$ are taken to be the same as that of the ρ^0 , which is parametrized as

$$\sigma_\rho = a + \frac{b}{\sqrt{E}}, \quad (3.29)$$

where $E \approx \nu$ is the energy of the meson. Note that this behavior is rather different from the $\sigma_V \sim 1/M_V^2$ one of the other two models. When ν and Q^2 become large, the second term in 3.29 becomes small and shadowing is then approximately Q^2 independent, since at large Q^2 and fixed x the coherence length d is also roughly Q^2 independent. The amplitudes of the transitions $V_i \rightarrow V_j$ are adjusted so as to obtain Q^2 independence at large Q^2 and a good fit to the nucleon electroproduction data. Figure 3.24 shows the comparison with the NMC results [38] (the predictions correspond to the measured Q^2 values). The agreement with the experimental data is good for $x > 0.01$; at lower values of x the theoretical curve is above the data.

None of the described VMD models makes explicit predictions for the very small x region ($x < 10^{-3}$) covered by the FNAL E665 results [41], where the cross section ratio Xe/D appears to saturate. A description of this region is however possible. In a simple VMD model with ρ , ω and ϕ contributions only [182] (see also [177]) one obtains saturation for $\nu > 6$ GeV (at $Q^2 = 0$), with a cross section ratio σ^{Xe}/σ^D approximately equal to 0.65. For $Q^2 = 0.01$ and $Q^2 = 0.1$ the predicted ratio becomes 0.67 and 0.84, respectively [209]; the former value is in good agreement with the E665 results.

3.2.2 Partonic Models in the Breit Frame: the Recombination Approach

The first discussion of shadowing in a partonic approach is the one of N.N. Nikolaev and V.I. Zakharov [195]. Parton-nucleus scattering is considered in the Breit frame, where the energy of the virtual photon is zero and the three-momentum of the scattered parton is exactly reversed by the collision. A parton carrying a fraction x of the nucleon's momentum P_N can only be localized longitudinally to within a distance $\Delta z \sim 1/(xP_N)$ (the "longitudinal dimension" of the parton), as can be seen by using the uncertainty principle. On the other hand the nucleons are separated by a distance

$$\Delta z_N \sim 2R_N$$

(with the nucleon radius $R_N \approx 1$ fm) in the laboratory frame or

$$\Delta z_N \sim 2R_N \frac{M}{P_N} \quad (3.30)$$

in the Breit frame. So for $x < x_N = 1/(2R_N M) \approx 0.1$ partons belonging to different nucleons at the same impact parameter start to overlap spatially. For x values smaller than $x_A = 1/(2R_A M) \approx x_N A^{-1/3}$ (R_A = the nuclear radius), Δz exceeds the Lorentz contracted nuclear diameter and partons from all nucleons (at the given impact parameter) overlap. Overlapping partons can interact and fuse, thereby reducing the parton density at small x .

The total momentum carried by partons is not changed as a result of the fusion but is simply redistributed to different x regions. The depletion of low momentum partons due to the fusion process thus leads to an enhancement in the density of higher momentum partons, expected to be maximum around $x \approx x_N$.

Since in DIS the cross section is proportional to the parton distribution functions, the depletion of low x partons causes a depletion in the nuclear structure functions which starts somewhere below x_N and becomes maximum at $x \approx x_A$. Similarly the enhancement in the parton density at $x \approx x_N$ determines an increase – “anti-shadowing” – of the bound nucleon structure functions with respect to the free nucleon ones. The prediction of anti-shadowing is a feature characteristic of the parton recombination approach. Also characteristic of this treatment, and in contrast to the VMD models, is the fact that shadowing and anti-shadowing do not disappear with increasing Q^2 , but are rather functions of x alone (they “scale”). The qualitative trend of the ratio according to Nikolaev and Zakharov is shown in fig. 3.25.

More recently Müller and Qiu [196] approached parton fusion in a perturbative QCD framework and studied the consequences of gluon recombination. They derived modified, non-linear evolution equations which include gluon recombination effects. Such effects turn out to be significant only if evolution is started below $\approx 2 \text{ GeV}^2$. They concluded that the main consequence of gluon recombination is a depletion of the gluon distribution in nuclei at small x which reflects into a depletion of the sea quark density; the direct shadowing of sea quarks themselves appears instead to be a smaller effect.

By means of these modified Altarelli-Parisi equations, Qiu [197] focused on the Q^2 evolution of shadowing. The nuclear parton distribution functions used as an input are a product of a phenomenological shadowing function R_s and a parametrization of the parton distributions in the nucleon. The explicit form for R_s , often referred to in the later literature, is:

$$R_s(x, Q_0^2, A) = \begin{cases} 1 & x_N < x < 1 \\ 1 - K(A^{1/3} - 1) \frac{\Delta V_A}{V_A} & x_A < x < x_N \\ 1 - K(A^{1/3} - 1) & 0 < x < x_A, \end{cases} \quad (3.31)$$

where the constant factor K is different for sea quark and gluon distributions and must be determined by comparison with the experimental data. The quantity

$$\frac{\Delta V_A}{V_A} = \frac{1/x - 1/x_N}{1/x_A - 1/x_N} \quad (3.32)$$

is a measure of parton overlap which is zero at $x = x_N$ and maximum at $x = x_A$. The term $(A^{1/3} - 1)$ gives the average number of shadowed nucleons: at very small x the photon cannot distinguish between partons of different nucleons at the same impact parameter; for large enough A , say $A > 8$, the effective number of nucleons thus equals the number of surface nucleons, $\approx R_A^2/R_N^2 = A^{2/3}$, and the number of “shadowed” nucleons (per effective nucleon) is $n_s = (A - A^{2/3})/A^{2/3} = A^{1/3} - 1$. Note that expression 3.31 just parametrizes the x dependence

of shadowing, but does not predict it. The main conclusion of Qiu is that the Q^2 dependence of shadowing is very weak, in agreement with the original prediction of Nikolaev and Zakharov.

After the publication of the EMC NA2' and NA28 results, Berger and Qiu [198] elaborated on this approach, concentrating on the A dependence of the point where shadowing sets on. These authors observed that the average longitudinal distance between one parton from a nucleon and another from a neighboring nucleon is larger for nucleons on the surface of the nucleus than for nucleons well inside the nucleus. As a result of these surface versus volume effects, the onset of shadowing is seen to increase with A :

$$\hat{x}_N = \frac{1}{MR_N \left[1 + 3 \frac{2}{A^{1/3}} - 3 \left(\frac{2}{A^{1/3}} \right)^2 + \left(\frac{2}{A^{1/3}} \right)^3 \right]}, \quad (3.33)$$

where the new symbol for the onset of shadowing is meant to emphasize the fact that $\hat{x}_N \neq 1/(2R_N M)$. The value of $\hat{x}_N$ rises from $\hat{x}_N = 0.08$ for ${}^4\text{He}$ to 0.1 for ${}^{12}\text{C}$ and 0.107 for ${}^{64}\text{Cu}$. While this trend is in qualitative agreement with the data, it is not entirely clear however what the actual meaning of $\hat{x}_N$ is (whether for instance it should be identified with the x value corresponding to the maximum of the structure function ratio, in the anti-shadowing region, or with the point where the ratio is unity or, more probably, with some other point in between).

While retaining the partonic mechanism of Qiu [197], Close and Roberts [199] later attempted to justify the A dependence of the onset of shadowing with arguments based on the Q^2 rescaling model which, as we have seen, may be used to describe the structure function ratio for x values between 0.2 and 0.6. In this model the nuclear structure functions at a specific Q^2 are given by the free nucleon structure functions at a larger value of Q^2 :

$$F_2^A(x, Q^2) = F_2^N(x, \xi(Q^2)Q^2), \quad (3.34)$$

where the A dependence of the parameter ξ is controlled by the average nucleon density. If the density is high there is a sizable probability for the nucleons to overlap, thereby modifying the confinement scale of quarks. As the density of nucleons grows with increasing A , the distance between neighboring nucleons decreases; so does the distance between partons, thus making it more likely for them to fuse. As a consequence the onset of shadowing $\hat{x}_N$ grows with A . With the overlap probabilities evaluated in [120], the results found for carbon, copper and tin are $\hat{x}_N = 0.121, 0.130$ and 0.135 , respectively. The trend is similar to that found by Berger and Qiu [198] and also seen in the data. Again it is difficult to identify the actual meaning of $\hat{x}_N$ in terms of measured quantities, but the authors claim that at least its variation with A should be analogous to that of the point where the ratio equals unity. In the EMC NA2' data [35, 36] this cross-over point moves by ≈ 0.06 in going from carbon to tin, as opposed to the predicted variation of $\hat{x}_N$, $\Delta\hat{x}_N = 0.014$. The discrepancy motivated Close and Roberts to propose that the value of x_A , where shadowing saturates, is A independent, unlike what was previously assumed. The A independence of x_A means that shadowing saturates already when Δz exceeds only few nucleon dimensions. The assumption is made that this occurs when at least 12 nucleons are involved. In other words x_A is taken to be a constant for $A > 12$; its value is set to $x_A = 1/(2MR_A)$ with R_A the radius of the carbon nucleus, resulting in $x_A \approx 0.043$. This assumption is combined with the A dependence of $\hat{x}_N$ found with the Q^2 rescaling hypothesis. The nuclear structure functions are then computed by applying Q^2 rescaling to the parton distributions at all x ; the sea quark distribution in the nucleus is

suppressed by the shadowing function 3.31, with $K \approx 0.1$ and the values of $x_N = \hat{x}_N$ and x_A just discussed. Figure 3.26 shows that the predictions thus obtained are in fair agreement with the EMC NA2' data [35].

Detailed predictions based on the partonic approaches by Qiu [197], Berger and Qiu [198] and by Close and Roberts [199] were obtained in [200] for the kinematic range covered by the NMC results.

Berger and Qiu [198] do not make explicit predictions on the x dependence of shadowing, but rather restrict themselves to an estimate of the x value where shadowing sets on. These predictions were combined in [200] with the ansatz 3.31 proposed by Qiu [197], which gives the amount by which the sea quark distribution in a bound nucleon is suppressed with respect to the free nucleon case. In this equation the quantity x_N (where shadowing should start) is assumed to have the behavior given by expression 3.33, while the point where shadowing saturates is taken to vary with A as $x_A = 1/(2MR_A) \approx 0.1A^{-1/3}$ (with R_A the nuclear radius).

In order to obtain a prediction covering the large x domain as well, in [200] both the Q^2 rescaling prescription [120] (equation 3.12) and the x rescaling one [92] (equation 3.3) were considered.

Figures 3.27 and 3.28 show the comparison of the model predictions (continuous curves) with the NMC results [38], for the Q^2 and x rescaling cases, respectively. While the predictions for C/D are in reasonable agreement with the data when Q^2 rescaling is used, in all other cases the model calculations fail to reproduce the experimental results. The predictions for Ca/D change slightly if the values of x_N and x_A in equation 3.31 are replaced by those suggested by Close and Roberts [199] (dashed lines in figs. 3.27 and 3.28), who evaluated x_N using the nucleon-nucleon overlap probabilities of the Q^2 rescaling model [120] and x_A assuming that saturation occurs at $x = 1/(2MR_A)$, where R_A is now the radius of the carbon nucleus, yielding $x_A \approx 0.04$. Also in this case the comparison with the data is rather poor.

In the papers just discussed ([196]-[199]) the x dependence of shadowing was essentially assumed by using phenomenological parametrizations. Close, Qiu and Roberts [201] attempted a semi-quantitative derivation of the shape of shadowing based on the parton fusion picture corroborated by QCD arguments. The results can be summarized as follows. Perturbative QCD studies how parton distributions evolve as a consequence of quarks radiating gluons, gluons converting into $q\bar{q}$ pairs and gluons splitting into gluon pairs. Quarks and gluons can of course also fuse, but these diagrams are generally not studied individually; rather their effects are subsumed into the input parton distributions. Fusion diagrams are instead explicitly considered in this paper and the equivalent of the parton splitting functions for the fusion processes $qg \rightarrow q$, $gg \rightarrow g$ and $q\bar{q} \rightarrow g$ ("parton fusion functions") are evaluated.

The results are applied to the case of nuclear structure functions, starting from the situation in which gluons and sea quarks leak from a nucleon into a neighboring one and recombine with the quarks and gluons of that nucleon. Leakage of gluons and leakage of sea quarks are studied separately. Figure 3.29 shows the resulting effect on the gluon distribution $xG(x)$ for Fe at $Q^2 = 5 \text{ GeV}^2$; it is qualitatively consistent with the NMC results [56]. The effect mainly arises from $gg \rightarrow g$ processes, rather than from $q\bar{q} \rightarrow g$ or $qg \rightarrow q$ ones. It turns out that only about 0.1% of the nuclear momentum is transferred from quarks to gluons: momentum conservation is thus assured almost separately within the individual parton species.

The distinction is made between "initial state recombination", with which this paper is concerned (fig. 3.30(b)), and the so-called "radiative recombination" (fig. 3.30(d)). Initial state recombination dominates at small Q^2 ; for higher Q^2 the net effect is a mix of the two types of recombination. The Q^2 dependence of radiative recombination was the subject of the

papers by Qiu [197] and Müller and Qiu [196], presented above.

In principle one should thus first compute the effects of initial state recombination to obtain the x dependence of the parton distributions at some low scale Q_0^2 ; these should then be fed into the evolution equations in order to obtain the parton distributions at larger Q^2 .

The picture that emerges is consistent with the known experimental facts about shadowing. A quantitative evaluation of the shape of shadowing is however precluded by the ignorance of the value of Q_0^2 at which the initial recombination mechanism is dominant over the radiative one and from which evolution begins.

A somewhat different approach is the one developed by Zhu and collaborators [202]-[207] within the framework of the constituent quark model [210, 211]. In this approach the nucleon is viewed as composed of three constituent quarks or “valons”. Valons are themselves composite objects, made up of the familiar quarks and gluons; the nucleon structure functions are thus the convolution of the constituent quark distribution G_N^c with the distribution F_N^c of the quarks and gluons in a constituent quark:

$$F_N(x, Q^2) = \sum_i \int_x^1 dy G_{N_i}^c(y) F_{N_i}^c(x/y, Q^2), \quad (3.35)$$

where the sum is over the three valons. The constituent quark distribution G_N^c is normalized to unity and each constituent quark is assumed to carry, on the average, $1/3$ of the longitudinal momentum of the nucleon.

The shape of G_N^c is determined by the interactions among the constituent quarks and is independent of the probe. The parton distribution F_N^c in a constituent quark depends instead on the Q^2 of the probe. At small Q^2 the virtual photon “sees” a structureless constituent quark (the valence quark, effectively); as Q^2 increases the structure of the constituent quark evolves according to QCD and the photon sees the usual sea quarks.

The nuclear medium influences both the structure functions F^c of the constituent quarks and the shape of their distribution functions G^c , but does not affect the normalization of G_N^c nor the momentum carried on the average by each constituent quark.

Let us begin with F^c . The main point here is that the gluons and the sea quarks of the three constituent quarks recombine with those of other bound nucleons so that the sea distribution is hardened. This is the usual recombination approach and gives rise to shadowing and anti-shadowing.

More in detail, several mechanisms are at work. In the QPM with QCD, sea quarks are produced by the splitting process $g \rightarrow q\bar{q}$, while gluons essentially originate from the QCD bremsstrahlung of valence quarks. Shadowing of the sea quarks goes predominantly via $q\bar{q} \rightarrow g$ recombination and via the gluon-gluon fusion process $gg \rightarrow g$. Shadowing of the gluons proceeds instead only via gg fusion, since shadowing of the valence distribution is assumed to be negligible. It is thus expected that shadowing of the sea distribution is stronger than that of the gluons. Similar conclusions were also reached in [196, 197].

With increasing Q^2 , the rate of $q\bar{q}$ and gg recombination processes increases since such processes are proportional to the q , $\bar{q}$ and g densities squared, which increase with Q^2 . However, the increase of the gluon density determines an increase of the $q\bar{q}$ production rate, which compensates the enhanced rate of recombination; since gluons are less shadowed than sea quarks, there is in fact some over-compensation and a logarithmic decrease of the size of sea quark shadowing is expected. Remarkably also anti-shadowing is predicted to decrease and eventually disappear at large Q^2 values ($\sim 40 \text{ GeV}^2$). Notice that the $q\bar{q} \rightarrow g$ process transfers

momentum from the quarks to the gluons and that the net momentum fraction transferred rises with Q^2 . This conclusion is at variance with that reached by Close, Qiu and Roberts [201].

The x dependence of shadowing is obtained with a prescription similar to that of Qiu (cf. eq. 3.31): at $x = 0$ one has

$$F_A^c(Q^2) = F_N^c(Q^2)(1 - n_s K), \quad (3.36)$$

with n_s the number of shadowed nucleons (per surface nucleon), $n_s = (A^{1/3} - 1)$ for large A . For larger values of x the following ansatz is made:

$$F_A^c(x) = \begin{cases} \alpha(1-x)^\beta & 0 < x < x_N \\ [\alpha(1-x_N)^\beta - F^c(x_N)] f(x) + F^c(x) & x_N < x < 1, \end{cases} \quad (3.37)$$

where $F^c(x)$ is the distribution in the absence of shadowing and anti-shadowing and the function $f(x)$ gives the magnitude of anti-shadowing.

We come now to the effects of the nuclear medium on G^c . The constituent quarks are assumed to be bound by a harmonic oscillator potential and the effect of the nuclear medium is expressed in terms of a weakening of the spring constant. This picture is not unique and an equivalent description can be obtained by increasing the constituent quark confinement size: for instance $R_{F_c}/R_D - 1 = 7.5\%$. The effect is in both cases a softening of the constituent quark distribution G^c ; the degree of softening is related to the nuclear density. In practice then, like in many of the models discussed so far, the description of the shadowing mechanism is combined with a model of the EMC effect at larger values of x .

Figure 3.31 compares the NMC results for helium, carbon and calcium [38] to the predictions of this model. In the figure the curves for $Q^2 = 1, 10$ and 40 GeV^2 are given.

The predictions for the three nuclei are in excellent agreement with the measurements. In particular one should remark that the small x points, which have small Q^2 ($< 2 \text{ GeV}^2$ for $x < 0.01$) are closest to the prediction for 1 GeV^2 . For $x \approx 0.1$, where Q^2 is about 10 GeV^2 , the curve for 10 GeV^2 shows the best agreement with the data. The expected Q^2 dependence seems thus to be confirmed. The predicted Q^2 slope for calcium at $x \approx 0.1$, $b \approx -0.01$, is consistent with the measured one.

Within the framework of this model, the integrals of the structure function differences were also computed [207]. The predicted values, for the range $0 < x < 0.8$, are -0.87×10^{-3} , -1.10×10^{-3} and -1.90×10^{-3} for He/D, C/D and Ca/D, respectively, compatible with the results of table 2.2, which however cover the range $0.0035 < x < 0.80$ only³.

Figure 3.32 shows the predictions of the model for the C/Li and Ca/C compared with the NMC results [39]. The agreement is good, except for the Ca/C ratio where the prediction is lower than the data at small x .

The model also describes well the E772 Drell-Yan results [72], as fig. 3.33 shows. Finally fig. 3.34 indicates that the enhancement in the Sn gluon distribution with respect to that in C observed by NMC [56] is consistent with the model predictions [208].

Before closing the discussion about this model, it should be noted [209] that it is not possible to extend its predictions to the region covered by the FNAL E665 data [41], the Q^2 of the data being too small for any partonic approach to be applicable. Actually, in the Q^2 region in which the model is valid, no saturation of shadowing is predicted.

³The contribution from the region $0 < x < 0.0035$ is expected to be rather small, and approximately amounts to -0.1×10^{-3} , -0.2×10^{-3} and -0.4×10^{-3} for He, C and Ca, respectively [42].

3.2.3 Partonic Models in the Target Frame: Pomeron Exchange

Castorina and Donnachie [212, 213] also argued that shadowing can be explained at the partonic level, but adopted a different framework. Their starting point is that the small x behavior of the structure function F_2 is controlled by pomeron exchange (see fig. 3.35). In this approach the virtual photon scatters off a pomeron emitted by the nucleon, thus probing the parton content of the pomeron rather than that of the nucleon itself. The authors suggest that in a nuclear environment the effective coupling β_q of the pomeron to a quark is suppressed because of nucleon overlap, i.e., with reference to fig. 3.35, $\tilde{\beta}_q < \beta_q$. This overlap effectively reduces the volume occupied by each nucleon and decreases the mean distance between the quarks of different nucleons. The hypothesis that the spatial separation between quarks (i.e. the radius of the hadron) determines the strength of the quark-pomeron coupling is consistent with the observation that the total hadron-hadron cross section is proportional to the squared radii of the interacting hadrons [214]. As the total cross section is also proportional to the effective quark-pomeron coupling, its variation can be interpreted in terms of a dependence of the pomeron coupling on the hadron radius [215, 216]. Quantitatively the structure functions at small x are then expected to be suppressed in a bound nucleon with respect to a free one by a factor proportional to the average nucleon-nucleon distance squared. This treatment of the low x behavior of nuclear structure functions is coupled with a description of the large x ($0.3 < x < 0.6$) behavior of the EMC effect based on the x rescaling model of Akulinichev *et al.* [92]. The description of the enhancement at intermediate x is left to the interplay between shadowing and x rescaling.

The resulting prediction for the ratio of bound to free nucleon structure functions is shown in fig. 3.36 together with the NMC results [38]. The model reproduces the data in the C/D case for $x < 0.01$ but underestimates them for larger values of x . The situation is worse for calcium where the expected shadowing signal and the experimental points cross each other at $x \approx 0.15$ but are otherwise in significant disagreement.

In a later paper [213] Castorina and Donnachie adjusted the normalization of the sea quark distribution in order to achieve better agreement with the EMC NA28 [37] and NA2' [35] data. The new predictions are shown in fig. 3.36 as dashed curves. The agreement with the NMC results is now better, especially for $x > 0.04$. At smaller x the model prediction is however systematically lower than the data, by up to 5-6% in carbon.

Preparata and Ratcliffe [217] also base their interpretation of shadowing on the dominance of pomeron exchange in the small x region. As a consequence of the diffractive nature of the interaction the nucleon on which the scattering has occurred is left essentially intact. If the nucleon is bound in a nucleus a large fraction of the phase space volume is not available to the scattered nucleon, because of the Pauli exclusion principle. The resulting predictions are compared with the NMC results for C/D and Ca/D in fig. 3.37.

3.2.4 Pomeron Exchange and VMD Combined

Kwieciński and Badelek [218] proposed a model which combines both the rescattering of low mass vector mesons and the partonic mechanism of shadowing in order to describe nuclear structure functions in the region of $0.001 < x < 0.1$ and $0.1 < Q^2 < 10 \text{ GeV}^2$. In their model, the structure function F_2^A of a bound nucleon is related to the free nucleon structure function by the following expression:

$$AF_2^A = AF_2^N - \Delta F_2^{(v)} - \Delta F_2^{(p)}; \quad (3.38)$$

$\Delta F_2^{(v)}$ and $\Delta F_2^{(p)}$ are the vector meson and partonic contributions to shadowing, respectively.

- The vector meson contribution $\Delta F_2^{(v)}$ is obtained as usual by means of relation 3.23; the sum is restricted to the ρ^0 , ω and ϕ mesons only. The interactions of vector mesons with the nucleus are treated in the Glauber approximation and the nucleus is taken to be a sphere of uniform density and radius $R = r_0 A^{1/3}$.
- The partonic term $\Delta F_2^{(p)}$ is discussed in the pomeron exchange language. Figure 3.38 shows the forward elastic photon-nucleon (Compton) scattering diagram; at small x the blob in the figure corresponds to the photon-pomeron interaction shown in fig. 3.39. When the target is a nucleus rather than a nucleon, the multiple scattering terms of fig. 3.40 must be considered, in which the double interaction diagram of fig. 3.41(a) dominates. At small x this process is controlled by the double pomeron exchange of fig. 3.41(b). The pomeron structure function in this figure has two major components: the triple pomeron (fig. 3.41(c)) and the quark box diagram contribution (fig. 3.41(d)).

In the region $0.75 < Q^2 < 1.5 \text{ GeV}^2$ shadowing turns out to be practically independent of Q^2 . The pomeron exchange contribution is instrumental in this respect, since the contributions due to the rescattering of vector mesons decrease significantly over this region of Q^2 .

Kwieciński [219] later elaborated on the same approach in an attempt to establish a quantitative relationship between the shadowing corrections to nuclear parton distributions and the partonic content of the pomeron. Starting from the hypothesis that the pomeron is essentially composed of gluons at $Q^2 \approx 2 \text{ GeV}^2$ (in the sense that gluons carry all the pomeron's momentum) and assuming specific momentum distributions for these gluons, the amount of shadowing in nuclei can be estimated. In particular, shadowing of the nuclear gluon distribution is shown to be substantial: at $x = 0.01$ in iron it amounts to as much as 50%.

Figure 3.42 shows the NMC results for the carbon and calcium to deuterium structure function ratios [38] and the predictions of the model by Kwieciński and Badelek [218]. The upper and lower continuous curves correspond to $r_0 = 1.5$ and 1.25 fm , respectively, in the formula $R = r_0 A^{1/3}$ for the nuclear radius. Experimental data and theoretical predictions disagree, especially at high x ; the disagreement is more conspicuous in the carbon case. In fig. 3.42 we also show the prediction by Kwieciński [219] for the iron to deuterium structure functions ratio (dashed line). The curve was computed for $Q^2 = 10 \text{ GeV}^2$.

While in the previous model VMD and pomeron exchange contributions were taken as distinct and summed, in those that we are about to discuss they become two aspects of the same process.

In the approach of Brodsky and Lu [220] the virtual photon is assumed to convert to a $q\bar{q}$ pair at a distance before the target proportional to $\sim 1/x$ (cf. expression 3.22), in the laboratory frame. In the Bjorken limit and in an appropriate gauge the final state interactions of the quark can be neglected in the valence approximation and effectively only the anti-quark interacts. The anti-quark-nucleus scattering amplitude is obtained from the anti-quark-nucleon one by means of Glauber's multiple scattering series. At very small x the $\bar{q}N$ amplitude is taken to be dominated by pomeron exchange and is mainly imaginary; as we emphasized earlier, this

gives rise to shadowing. In order to reproduce the enhancement at larger values of x , the authors argue that the real part of the amplitude becomes important in this region; in the Regge language this amounts to describing the scattering process in terms of a superposition of reggeons.

Figure 3.43 shows the comparison of the model predictions with NMC results [38]. While the agreement is satisfactory for the carbon to deuterium ratio, it is somewhat worse for calcium. Note however that the prediction is actually for Cu rather than for Ca.

Frankfurt and Strikman [221, 222] examined the nuclear effects on the quark distributions from a similar starting point. At small x , in the laboratory frame, the virtual photon-nucleus interaction proceeds in two stages:

1. the virtual photon dissociates into a quark-anti-quark pair;
2. upon arrival on the nucleus, this virtual system interacts with the nucleons giving rise to the hadronic final state.

The interaction cross section of the $q\bar{q}$ pair is taken to be proportional to its transverse dimension squared, approximately given by $1/k_t^2$, with k_t the relative transverse momentum of the pair. Pairs with small k_t have large transverse size, comparable with that of a hadron, and indeed interact in a way resembling a hadron projectile. Large k_t pairs have transverse dimensions smaller than those of an ordinary hadron and their cross section with the nucleus is simply proportional to the number of nucleons. Small k_t and large k_t pairs are often called “asymmetric” and “symmetric”, respectively, with reference to the partition of the photon’s momentum between the two quarks: when one of the quarks carries nearly all the momentum of the photon, the relative k_t is small; k_t is instead maximized by pairs in which each quark carries the same fraction of the photon’s momentum.

The decrease of the opacity of nuclear matter with decreasing transverse dimensions of the $q\bar{q}$ pairs is often referred to as “color-transparency”. The distinction between low k_t asymmetric and large k_t symmetric pairs is a reformulation of the aligned jet model of Bjorken and Kogut [223], in which the nucleus was supposed to become less and less opaque to hadronic jets with increasing k_t of the jet. The idea of discriminating between symmetric and asymmetric pairs in order to give a description of shadowing was first proposed by Nikolaev and Zakharov [195].

Sea and valence quark distributions are treated separately.

- Sea quarks. For $x < x_A$ ($x_A = 1/(2MR_A)$) the cross section $\sigma_{\gamma A}$ reads:

$$\begin{aligned} \sigma_{\gamma A} = & \frac{\alpha}{3\pi} \int_{M_0^2}^{\infty} \frac{dM_V^2 R_{e^+e^-}(M_V^2) M_V^2}{(Q^2 + M_V^2)^2} \left[\frac{3k_{t0}^2}{M_V^2} \right] \sigma_{VA} \\ & + \lambda A \sigma_{\gamma N}(x, Q^2). \end{aligned} \quad (3.39)$$

The first term in 3.39 has the usual GVMD-like form and is the contribution of small k_t pairs, which behave like hadrons, in fact vector mesons of mass M_V . The second term is the contribution of high k_t pairs, which interact with non-hadronic, pointlike cross section.

Here $M_0^2 > m_\rho^2$ is the scale characteristic of soft hadronic processes and k_{t0} is a typical transverse momentum of such processes. The vector meson-nucleus cross section is σ_{VA}

and the nuclear radius is indicated as before by R_A . The quantity $R_{e^+e^-}(M_V^2)$ is the ratio between $\sigma(e^+e^- \rightarrow \text{hadrons})$ and $\sigma(e^+e^- \rightarrow \mu^+\mu^-)$. The main difference in the integrand in 3.39 with respect to that in formula 3.28 is the presence of the factor $3k_{t0}^2/M_V^2$, proportional to the phase space occupied by large transverse size pairs. This extra power of M_V^2 in the denominator of the integrand allows the $1/Q^2$ scaling behavior of $\sigma_{\gamma A}$ to be recovered with vector meson-nucleus (or -nucleon) cross sections which do not depend strongly on M_V^2 , but rather are constant, as intuitively expected (cf. the discussion in section 3.2).

For $Q^2 \gg m_\rho^2$ and $x \ll x_A$, equation 3.39 predicts

$$\frac{A_{eff}}{A} = (1 - \lambda) \frac{\sigma_{\rho A}^{tot}}{\sigma_{\rho N}^{tot}} + \lambda \simeq 1.3A^{-0.18}(1 - \lambda) + \lambda, \quad (3.40)$$

provided that $A > 10$. The parameter λ can be extracted from the experimental data and is found to be $\lambda \approx 0.2-0.3$.

The x dependence of shadowing is computed, as usual, by means of Glauber's multiple scattering series. The region where shadowing starts to saturate moves to smaller x with increasing A ; the dependence of this value of x on A is estimated to be slower than $\sim A^{-1/3}$.

- Valence quarks. Shadowing of the valence quark distribution V_A (i.e. of $\sigma_{W^+A} - \sigma_{W^-A}$) is calculated by applying Glauber's series to describe the interaction of the soft hadron component of the virtual bosons W^+ and W^- . At very small x ($\ll x_A$) and $Q^2 \approx 1$ GeV² the valence distribution is calculated to be also shadowed by as much as 0.6 for $A \approx 60$ and 0.5 for $A \approx 200$.

In order to explore the region of higher values of x , up to $x \approx 0.1$, Frankfurt and Strikman use momentum conservation

$$\int_0^A \frac{1}{A} [V_A + S_A + G_A] \xi_A d\xi_A = \int_0^1 [V_N + S_N + G_N] x dx \quad (3.41)$$

and baryon number conservation

$$\int_0^A \frac{1}{A} V_A d\xi_A = \int_0^1 V_N dx, \quad (3.42)$$

where $\xi_A = Ax$ and V, G and S are the valence quark, gluon and sea quark distributions, respectively.

Since the valence distribution is shadowed at small x , baryon number conservation implies that valence is enhanced at higher x . Likewise shadowing of the sea and momentum conservation imply an enhancement of the sea and/or of the gluon distribution at higher x .

Starting from the experimental results on the integral of $F_2^A - F_2^D$, Frankfurt, Strikman and Liuti [224, 225] estimated the ratio of momentum fractions carried by gluons in Ca and D and found it to be larger than unity by a few per cent. Combining the conclusions reached above on the shadowing of sea and valence quark distributions with these findings, they derived the

shape of the parton distributions in nuclei. All parton distributions turn out to be shadowed at small x , while at larger x only the valence quark and gluon distributions are enhanced.

The dynamical mechanism suggested to explain this pattern is once again parton fusion, when DIS is viewed in the Breit frame. Here shadowing of the valence distribution is the result of the absorption of gluons by valence quarks. For the sea two processes compete: $q\bar{q} \rightarrow g$, which depletes the sea quark distribution and $qg \rightarrow q$, which enhances it.

The comparison with the NMC results [38] is shown in fig. 3.44 for the Ca/D ratio. The model calculation agrees well with the data in the anti-shadowing region; it is lower by up to 4-5% for smaller values of x . The Q^2 dependence predicted by the model is consistent with the experimental results, with logarithmic slopes $d(F_2^A/F_2^D)/d(\ln Q^2)$ positive up to $x \approx 0.1$ and then turning negative for larger values of x .

The model predictions for the ratio of the sea quark distributions are in good agreement with the E772 results [72], as fig. 3.45 shows; from the figure it is clear that the expected Q^2 dependence for the Drell-Yan ratio is small.

Finally the prediction of an enhancement of the gluon distribution around $x \approx 0.1$ is consistent with the NMC results [56]; we recall however that the NMC data are for the Sn/C ratio.

A further contribution to the discussion on shadowing came recently in two papers by N.N. Nikolaev and B.G. Zakharov [226, 227], later complemented by a third one, in which shadowing and the EMC effect are discussed together [137] (see also section 3.1.3).

Once again shadowing is ascribed to the large transverse size, asymmetric, $q\bar{q}$ component of the virtual photon. Small size, symmetric pairs also contribute to shadowing, but their effect dies away quickly with increasing Q^2 .

This description is made quantitative by means of perturbative QCD. The interaction of the virtual photon with the nucleons is described in terms of pomeron exchange. Pomerons are effectively treated as gluon pairs. Figure 3.46 shows the diagrams involved, in which the gluons of the pomeron interact with the $q\bar{q}$ state coming from the "diffractive" dissociation of the photon. The diagrams in which the photon excites a $q\bar{q}g$ state are also included in the treatment ("triple pomeron vertex"). In the calculation of the small x nucleon structure function, based on these diagrams, use is made of the spatial wavefunction of the $q\bar{q}$ fluctuations. The extension to the nuclear case is then made via Glauber's multiple scattering approximation. The QCD analysis confirms the near-scaling of shadowing and predicts only a slow ($\sim 1/\ln Q^2$) evolution with increasing Q^2 . The authors point out the connection between this approach and the fusion one of [195]. For example the quarks in the $q\bar{q}$ loop of fig. 3.46 do not belong to either of the nucleons involved: in the language of [195], this is a typical fusion of the gluon of the first nucleon with the sea quark of the second nucleon, or vice versa.

In these papers it is also shown that shadowing should be different for the transverse and longitudinal absorption cross sections, σ_L and σ_T : shadowing of longitudinal photons is expected to disappear more quickly with Q^2 than that of transverse photons. A consequence of this is the A dependence of $R = \sigma_L/\sigma_T$, which should vary by up to 10-15% in going from hydrogen to lead.

Finally, an explicit prediction is made for shadowing of the deuteron structure function:

$$\left(\frac{A_{eff}}{A} - 1\right)_D \approx 0.25 \left(\frac{A_{eff}}{A} - 1\right)_{He} \quad (3.43)$$

All predictions given for F_2^A/F_2^N include such contribution. Note that this procedure is *not* generally adopted by other authors.

Figure 3.47 shows the comparison of the predictions presented in [137] with the NMC results [38]. These predictions differ from those in [226, 227], where the EMC effect was neglected. As fig. 3.47 shows in the He and Ca cases, the EMC effect calculation taken alone (and based on swelling of the nucleon, cf. section 3.1.3) predicts an enhancement of the structure function ratio. What is observed in the data is actually the combination of this enhancement and of the depletion due to shadowing.

The model prediction are in excellent agreement with the NMC data for He/D; they tend to slightly overestimate the C/D and Ca/D results in the $0.01 < x < 0.1$ region. The predicted Q^2 dependence is consistent with the measured one. The model is not applicable to the Q^2 covered by the FNAL E665 [41] results, in the region of $x < 10^{-3}$.

Figure 3.48 compares the model predictions with the E772 results [72]; the theoretical curves and the data are in fair agreement – the predictions are for $Q^2 = 4 \text{ GeV}^2$ while the data cover the region $16 < Q^2 < 81 \text{ GeV}^2$; the Q^2 dependence of this ratio is however expected to be small.

The model predicts an enhancement of the bound to free nucleon gluon distribution by up to 3% at $Q^2 = 4 \text{ GeV}^2$, for $x < 0.1-0.2$; this is qualitatively in agreement with the NMC results [56], which are however for the Sn/C ratio.

3.2.5 Conclusions of the Comparison with the Theoretical Models

The NMC data allow some discrimination between the models: it is now clear that some of these appeared to describe the EMC NA28 and NA2' results only because of the large error bars of the data. A few remarks are in order:

- Generalized vector meson dominance, twenty years after its original formulation, appears to offer still a viable description of shadowing. It is a little disturbing however that different versions of essentially the same model give predictions in disagreement with each other. Part of the problem is probably in the many adjustable parameters of the model (e.g. the vector meson-nucleon cross section, the spectrum of the high mass vector resonances etc.).
- As it has been emphasized by several authors (see e.g. [78, 198, 226, 227, 137]) a partonic description of shadowing based on the recombination approach is not necessarily in contradiction with GVMD, but rather expresses the same physical picture in the Breit frame; in the $q\bar{q}$ -pomeron picture VMD is actually recovered and “explained” at a more microscopic level. It is therefore not astonishing that partonic and GVMD models may be both consistent with the data.

The validity of partonic approaches may however extend to larger values of x to include the enhancement region. In a formulation based on parton fusion, the enhancement is a natural consequence of momentum conservation.

- Many models of shadowing neglect the possibility that the mechanisms leading to the EMC effect may be already at work at very small x . However, if this is the case, as suggested for instance in [226, 227, 137], the measured structure function ratios at small x are the sum of the shadowing and of the EMC effect contributions. The fact that some shadowing models agree with the data and some do not would become, in this case, largely accidental.

- Unfortunately very few of the models (none ?) that we have discussed are really predictive in the sense that they are proven incorrect if they do not describe the experimental data. In most cases the predictions can be “improved” just by careful tuning of the parameters. In a broad sense, for instance, the model by Zhu *et al.* [204]–[206], which gives an accurate description of the data, is a refined version of the one by Close and Roberts [199], which is not proven wrong, strictly speaking, by figs. 3.27 and 3.28; the poor comparisons simply indicate that some of the assumptions (e.g. that saturation occurs for nuclei with $A \geq 12$) or of the parametrizations adopted (e.g. the one describing the x dependence of shadowing) need to be revisited.

Unlike what happens in other branches of particle physics, the role of accurate data seems here that of a guide (sometimes a mere source of parameters) for theory, which for the moment appears to be far from being able to make predictions based on first principles.

We have so far concerned ourselves with the small x region ($x < 0.2$), where the recent data (NMC, E665, E772) have brought new information. Before concluding this section we want to briefly comment on the status of the comparison between theory and experiment the EMC effect region, for $x > 0.2$.

As we discussed at length, a variety of models is able to describe the EMC effect. Different physical approaches are used in order to make the quark distribution in a bound nucleon softer than that in a free one – which is the essence of the phenomenon. Most approaches give a reasonable description of the results for all the nuclei on which data are available. The x and A dependences of the effect are therefore not sufficient to discriminate among the existing models. If we include in the comparison also the Drell-Yan data [72] – to be taken however with some caution, as mentioned above – then the pion approach (and x rescaling ?) appears to be ruled out, as we discussed before. The statement contained in [72] that also other approaches (notably cluster models and Q^2 rescaling) are contradicted by these results has not met significant consensus: there are cluster models which can reproduce these Drell-Yan data (see e.g. [170] and fig. 3.18); furthermore the fact that Q^2 rescaling does not predict shadowing is well known, but does not automatically rule out the model in the EMC effect region.

Some more definite discrimination between the models of the EMC effect must probably await reliable data in the $x > 1$ region. If nuclear structure functions are proven to be non-zero there, then x and Q^2 rescaling models are ruled out, as well as all models based on nuclear binding or pions. In none of these models can a quark carry a fraction of the nucleon momentum significantly larger than unity. This is instead possible in a natural way in models based on quark deconfinement or on multi-quark clusters. Note that some attempts [10] made to describe the $x > 1$ region in terms of Fermi motion are in fact based on variants of cluster models.

Quark-cluster models might also offer, at least in principle, the possibility of a unified description of shadowing and the EMC effect. Indeed we observed earlier that the predictions of some of these models extend to x values well inside the shadowing region. Quark-cluster models in fact overlap to a certain extent with the parton recombination approach for shadowing. As observed in [168] the probability that a quark in a nucleus is part of a cluster may be taken as a crude measure of the amount of recombination. Unfortunately, the predictions of these models are in general rather vague in the shadowing region (cf. for example [165, 168] and fig. 3.15) and do not allow a stringent comparison with the data.

New light on the problem may be shed by splitting the A dependence of nuclear effects in structure functions into two parts, one depending linearly on A (“volume part”) and one on

$A^{2/3}$ ("surface part") [228, 229]. The study of the volume term is equivalent to that of nuclear effects in infinite nuclear matter, which may be easier to treat theoretically.

3.3 Summary of the Theoretical Situation

The detailed and extensive experimental information now available is confronted with a wealth of theoretical models which generally reproduce the qualitative features of the data but for the moment fail to be predictive at a quantitative level.

The following theoretical scenario is consistent with the data:

- At small x , the situation can be described in different ways:
 1. In a frame in which the target nucleons are at rest the virtual photon can be seen as fluctuating into a quark-anti-quark pair. The photon thereby acquires hadronic properties and can be effectively treated as a vector meson. The meson interacts strongly with the nucleons on the surface of the nucleus, which thus absorb a significant part of the incoming flux and cast a shadow on the inner nucleons. Of the $q\bar{q}$ pairs into which the photon fluctuates, only those with large transverse size do interact with hadronic cross sections; such pairs are a fraction of the total which decreases with increasing invariant mass of the pair.
 2. In the Breit frame, small momentum quarks and gluons, because of the uncertainty principle, spread over a distance comparable to the nucleon-nucleon separation. Quarks and gluons from different nucleons can overlap spatially and fuse, thus increasing the density of high momentum partons at the expense of that of lower momentum ones.

This point of view is attractive because it accounts, at the same time, for the existence of shadowing and of anti-shadowing. Also it implies that some of the momentum carried by quarks may be transferred to gluons via $q\bar{q} \rightarrow g$ processes.

It is nearly uncontested that sea quarks and gluons are the principal characters of shadowing, the role of valence quarks being small at low values of x . This is (perhaps – cf. the theoretical uncertainties mentioned in 2.3.10) confirmed by the fact that shadowing effects similar to those observed in DIS have been measured in Drell-Yan proton-nucleus dimuon production [72], sensitive to the target anti-quark distribution only.

The fact that the same Drell-Yan experiment found no clear evidence of anti-shadowing has instead led to contrasting conclusions. On the one hand, the absence of anti-shadowing in an experiment sensitive to anti-quarks suggests that anti-shadowing has to do with valence quarks (see e.g. [224, 78]). On the other hand, the disappearance of anti-shadowing in Drell-Yan can be a consequence of the measured Q^2 dependence in the enhancement region (cf. [204]-[206]).

- At larger values of x , $0.2 < x < 0.8$, the valence distribution appears to be softer in bound nucleons than in free ones. Many mechanisms have been proposed and most describe the results in a reasonable way. While the pion model appears to be ruled out by the FNAL E772 Drell-Yan data [72], for the others it will probably be necessary to wait for good data at $x > 1$ in order to operate a serious selection.

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$$|\gamma \text{ physical}\rangle = x \text{ wavy line } x + \sqrt{\alpha} \left(x \text{ wavy line } \text{hadron blob} + x \text{ wavy line } \text{e}^+ \text{e}^- \text{ loop} \right) + O(\alpha)$$

$$= |\gamma \text{ bare}\rangle + \sqrt{\alpha} |\text{hadrons}\rangle + \sqrt{\alpha} |e^+e^-\rangle + O(\alpha)$$

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