THE EMC EFFECT AND THE SWELLING OF NUCLEONS IN NUCLEI

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A B S T R A C T

We discuss the modifications of the electromagnetic polarizability and of the charge radius of a nucleon inside the nuclear medium. For a free nucleon, these quantities are largely influenced by the point cloud. In the nucleus the distortion of this cloud changes their value. We link these modifications to the quasi-deuteron part of the photoabsorption cross-section.

Invited Talk at the Heidelberg Symposium on
"Weak and Electromagnetic Interaction in Nuclei
1-5 July 1986

The EMC Effect and the Swelling of Nucleons in Nuclei

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The EMC effect has proved to be a stimulus for nuclear physics. Whatever the final word will be, it has triggered a number of investigations of great interest. Two schools of thought have emerged to interpret the EMC effect. The first one links it to some change in scale for a nucleon inside the nucleus, which entails the rescaling phenomenon [1]. This change in scale corresponds to a swelling of the nucleons, interpreted as a partial deconfinement of the valence quarks, a first step towards the total deconfinement of the quark-gluon plasma. This is a very exciting possibility, as it opens the way for a third generation of nuclear physics, after the traditional one and that of the meson degrees of freedom.

The second interpretation sees in the EMC effect a manifestation of phenomena which belong to the by now classical nuclear physics. In this case the influence of quantum chromodynamics (QCD) is only indirect as it is deeply buried in the physics of the nuclear forces or of the meson exchange theory. The most concise example is the description of the EMC effect in terms of the nuclear binding proposed by Garcia Canal et al. [2] and Akoulinitchev et al. [3]. These authors showed how the binding of the nucleons affects the structure function. The potential which binds the nucleons appears on the same footing as their kinetic energy which produces the Fermi motion correction previously considered. This effect can explain the softening of the valence quarks. Such a description asks no questions about the origin of the field which binds the nucleons. This binding is usually thought of as due to the meson exchange NN interaction but there are some attempts to link it more directly to QCD. The binding model does not reproduce the small x enhancement. In this sense, the pionic model is more complete [4]. Here the small x enhancement is attributed to the existence of a pion excess in nuclei and the softening of the valence quark to the loss of momentum of the nucleons to the benefit of the pions. The link between this model and the binding correction is contained [5] in the relation between the pion excess number per nucleon $\langle n_\pi \rangle$ and the expectation value $\langle V_\pi \rangle$ of the one pion exchange potential (OPEP):

$$\langle n_\pi \rangle = \langle V_\pi \rangle / \langle \omega_\pi \rangle,$$

where $\langle\omega_\pi\rangle$ is the average pion energy. The existence of a pion excess follows from the attractive nature of the OPEP potential. According to Ref. [5], pions provide most of the nuclear binding with $\langle V_\pi\rangle \approx -50$ MeV, which leads to a pion excess of $\langle n_\pi\rangle \approx 0.2$ per nucleon.

A recent work by Oka and Amado [6] seems to reconcile these two classes of interpretations. These authors show that the existence of a swelling follows from general grounds for a quantum system and is a consequence of the attractive nature of the binding potential. This step establishes a bridge between binding and swelling and justifies the existence of the swelling. So it seems that everybody should be happy, as all phenomena follow from each other. I am afraid that, in this talk, I will introduce some dissonance in this harmony by addressing the question of the swelling and of its interpretation.

The existence of a swelling has been looked for in low-energy nuclear physics. Are there any indications that a nucleon gets bigger in the nucleus? The answer is positive, there is an indication from inelastic electron scattering, from the Coulomb sum rule. Sum rules are particularly useful to disentangle nucleon properties from the nuclear ones. In the inclusive ee' scattering, the longitudinal response (charge) has been separated out [7] from the transverse one in the region of the quasi-elastic peak where the photon absorption occurs on one nucleon, a region which is well suited to measuring the nucleonic properties. The area under the longitudinal response should measure the total charge of the nucleus. This is the Coulomb sum rule. The puzzling result of these experiments is that $\approx 40\%$ of the charge is missing (Fig. 1). Noble was bold enough to suggest [8] that this is the signature of a swelling. Indeed, to obtain the charge, the nucleon form factor had been divided out, with the assumption that it is the same as that of a free nucleon. Noble [8] and later Shakin [9] proposed that this is not the case and that for a bound nucleon it drops faster than for a free nucleon, a clear signature of a swelling effect. The increase in the radius needed to recover the nuclear charge is $\Delta\langle r^2\rangle_p \approx 0.4$ fm², as compared to $\langle r^2\rangle_p = 0.65$ fm². This is the most natural interpretation of the missing-charge problem.

Fig. 1 The Coulomb sum rule for ^3He , ^{12}C , ^{40}Ca , ^{48}Ca , and ^{56}Fe . From Ref. [7].

The questions that I want to address now are: What is the physical interpretation of this swelling? What part of the nucleon structure is concerned? Does it justify the rescaling assumption? These questions lead to the simultaneous discussion of the root mean square (r.m.s.) charge radius and of the electromagnetic (e.m.) polarizability of a nucleon inside the nuclear medium. There are certain similarities between these quantities as the understanding of one helps in the understanding of the other. The polarizability is defined from the forward Compton amplitude, whose imaginary part is proportional to the total photoabsorption cross-section. It then becomes interesting to discuss our understanding of nuclear photoabsorption. Consider first the polarizability and the r.m.s. charge radius for a free nucleon and let us then discuss how these quantities are modified in the nucleus [10].

Free nucleon

The e.m. polarizability determines the leading t dependence of the forward and spin-averaged Compton amplitude $f(\omega)$

$$f(\omega) = f(0) + a\omega^2 + \dots, \quad (1)$$

where $a = \alpha + \beta$ is the sum of the electric and magnetic polarizabilities. This is a first similarity to the r.m.s. charge radius which fixes the t dependence of the electric form factor $G_E(t)$

$$G_E(t) = G_E(0) + (1/6) \langle r^2 \rangle t + \dots \quad (2)$$

In addition, both quantities a and $\langle r^2 \rangle$ arise from the internal structure of the nucleon, in fact from the most peripheral part of this structure which is dominated by the pion cloud. Its influence on the quantity a , for instance, can be gauged from the following dispersive integral

$$a_p = (1/2\pi^2) \int_{m_\pi}^{\infty} d\omega \sigma_p(\omega)/\omega^2, \quad (3)$$

where σ_p is the photoabsorption cross-section on the proton. In the cross-section itself the Δ resonance is very prominent but in the integral the denominator ω^2 strongly favours states near the pion threshold, i.e. πN states in the continuum [11]. Here the photoproduction cross-section has an electric (not a magnetic) dipole character. It occurs via the contact term or pion-pole one. The dominance of these contributions over that of the Δ explains why the electric dominates the magnetic polarizability in the quantity a ($\alpha \approx 10 \times 10^{-4} \text{ fm}^3$ versus $\beta \approx 4 \times 10^{-4} \text{ fm}^3$). A similar analysis based on dispersion relations gives the r.m.s. charge radius of the proton as [12]

$$\langle r^2 \rangle_p = (6/\pi) \int_{4m_\pi^2}^{\infty} dt \text{Im } G_E(t)/t^2, \quad (4)$$

where the spectral function $\text{Im } G_E(t)$ is dominated by two-pion states (Fig. 2). As in the previous case, the denominator t^2 favours the low-mass states near the two-pion threshold, a region where the two pions are nearly uncorrelated, rather than the region around $t = 30m_\pi^2$ where the ρ resonance occurs.

The importance of the pion cloud in the two structure constants has been confirmed by theoretical studies in the chiral bag model [13–15]. Oset et al. [14] find for the pion contribution to the r.m.s. charge radius: $\langle r^2 \rangle_{\text{pion}} \approx 0.5 \text{ fm}^2$. In the study of Theberge et al. [15] the cloud contribution is not as dominant, i.e. $\langle r^2 \rangle_{\text{pion}} = 0.2 \text{ fm}^2$, but it is by no means negligible.

In fact the analogy between the electric polarizability, which dominates the e.m. one, and the r.m.s. charge radius is more than formal. There is a link between these two quantities, which can be established as follows. The static electric polarizability is

$$\alpha = 2 \sum_n \frac{|\langle 0 | D | n \rangle|^2}{(E_n - E_0)} = (1/2\pi^2) \int d\omega \sigma_{\text{dip}}/\omega^2, \quad (5)$$

where D is the electric dipole operator which for a particle bound by a potential is $D = e\vec{r}$, and σ_{dip} is the unretarded electric dipole cross-section. On the other hand, the r.m.s. charge radius can be expressed with closure in terms of a sum over intermediate states

$$\langle r^2 \rangle = 3/(4\pi^2 e^2) \int d\omega \sigma_{\text{dip}}/\omega, \quad (6)$$

which provides the relation

$$\langle r^2 \rangle = (3/2e^2) \langle \omega_{\text{dip}} \rangle \alpha, \quad (7)$$

where $\langle \omega_{\text{dip}} \rangle$ is a typical E_1 energy.

Applied to the free proton, relation (7) gives an idea of the pion contribution to the proton charge radius. With the value of the intrinsic electric polarizability [16] of the proton $\bar{\alpha} = \alpha - (e^2/3M) \langle r^2 \rangle_p = 8.6 \times 10^{-4} \text{ fm}^3$ and a typical E_1 energy $\langle \omega_{\text{dip}} \rangle \approx 200 \text{ MeV}$, relevant for the πN continuum, the relation (7) gives $\langle r^2 \rangle_{\text{pion}} \approx 0.2 \text{ fm}^2$, which is of the right magnitude.

Fig. 2 The pion-cloud contribution to the r.m.s. charge radius of the proton

In the nucleus

The pion cloud is the longest range structure of the nucleon, extending in the nucleus in regions where other nucleons are present. Appreciable distortion of this cloud must then be expected in the nuclear medium [17]. This should have repercussions both on the polarizability and on the r.m.s. charge radius. As an illustration take a specific contribution to the proton e.m. polarizability, such as the one shown in Fig. 3. Here the pion is absorbed and re-emitted by the same nucleon. In the nucleus it can also be a pion in flight between two different nucleons, as shown in Fig. 4. In this case we obtain a meson-exchange correction to the nucleon e.m. polarizability. As an example, the evaluation of the meson-exchange correction, such as the one shown in Fig. 4, for a free Fermi gas, gives the Pauli correction to the e.m. polarizability. The overall Pauli correction results in a quenching of a_N by $\approx 20\%$, by no means a negligible effect [10]. According to the relation (7) a similar Pauli quenching should apply to the r.m.s. radius. We will now see that this is not the whole story but the message is already clear. A change in the r.m.s. charge radius of the nucleon inside the nucleus is more likely to reflect a distortion of the pion cloud, rather than a swelling of the quark core as is commonly thought. In order to see what further modifications occur, I will first describe the photoabsorption cross-section on nuclei, an ingredient which enters into the expression for the polarizability and for the r.m.s. charge radius.

Fig. 3 A contribution to the nucleon polarizability arising from the pion cloud

Fig. 4 A meson exchange correction to the nucleonic polarizability

Nuclear photoabsorption cross-section

This cross-section displays two strong resonances, the giant dipole resonance (GDR) at low energy and the Δ one at ≈ 300 MeV (Fig. 5). These two resonances are very different in nature. The GDR is a collective oscillation of the protons against the neutrons. Its position varies from nucleus to nucleus, with the law $\omega_{\text{GDR}} \text{ (MeV)} \approx 80 A^{-1/3}$. The Δ instead is an excitation of the individual nucleons. Its position is the

same for all nuclei. This difference is conceptually important for the understanding of the question of the e.m. polarization and of the r.m.s. charge radius on nucleons inside the nuclear medium. This appears in the investigation of the nuclear polarizability a_A described below. This quantity can be described by a dispersion relation similar to (3) with the integral starting at $\omega = 0$. We now want to separate out the contribution from the *individual nucleons*. For this we make the following observation. Suppose that in the cross-section σ_A there were only the two resonances, the GDR and the Δ , well separated, then the nucleonic contribution to a_A , $\tilde{a}_N$ could easily be singled out

$$\tilde{a}_N = (1/2\pi^2) \int_{m_\pi}^{\infty} [\sigma_A(\omega)/A]/\omega^2. \quad (8)$$

The comparison between the free polarizability a_N and that of a bound nucleon $\tilde{a}_N$ would amount to the comparison between the free cross-section and the nuclear one per nucleon above pion threshold, which are close. However, there is more than just these two resonances, which makes this separation more subtle. Above the GDR the photoabsorption cross-section does not vanish but it shows a plateau. This is the so-called quasi-deuteron (QD) region and the name arises from the belief that the absorption is a two-nucleon process which occurs on correlated pn pairs. The quasi-deuteron cross-section has been parametrized as a scaling of the deuteron one with the Levinger factor L' . In addition there is a damping factor $e^{-80/\omega}$, at low energy, which is due to the Pauli blocking that forces the emitted nucleons to be above the Fermi sea:

$$\sigma_{QD}(\omega) = L' (ZN/A) [62.5 (\omega - 2.2)/\omega^{5/2}] e^{-80/\omega}, \quad (9)$$

with σ in mb and ω in MeV.

Fig. 5 Experimental values of the nuclear photoabsorption cross-section per nucleon. From Ref. [18].

Microscopically, the quasi-deuteron absorption, up to $\omega \approx 140$ MeV, is a pionic process. The excitation of the 2p2h state occurs essentially via pion exchange, through the contact term (a) or the pionic current one (b) shown in Fig. 6. Some contribution also comes from the correlation term, one example of which is shown in Fig. 7. The important feature is that the quasi-deuteron absorption is an electric dipole mode. This does not hold on the nuclear scale but at the level of the pn pairs, which set the relevant scale for this absorption mechanism.

Coming back to the question of the nucleonic polarizability inside the nucleus, how should we consider the contribution to the nuclear polarizability of the quasi-deuteron cross-section? Does it pertain to the nucleus or the individual nucleon? To answer this, we observe that the quasi-deuteron cross-section has a *volume* character, in contradistinction to the GDR one which is a surface effect. It is thus normal to incorporate this part of the cross-section in the expression of the nucleon polarizability and to write $\tilde{a}_N$ as

$$\tilde{a}_N = (1/2\pi^2) \int_0^\infty d\omega [\sigma'_A(\omega)/A]/\omega^2, \quad (10)$$

where in σ'_A only the volume part of the cross-section has been retained: $\sigma'_A = \sigma_A - \sigma_{\text{GDR}}$. With the expression (10) and the parametrization (9) of σ_{QD} we find that the nucleonic polarizability in the nucleus is greatly enhanced, by a factor of ≈ 3 . The origin of this large enhancement is the existence of low-lying ($\omega \lesssim m_\pi$) 2p2h excitations which generate σ_{QD} . Notice that this enhancement concerns mostly the electric polarizability.

We would like now to understand this result in the language of the meson exchange picture, to translate the physical effects in the cross-section into the treatment with virtual pions. In order to illustrate the connection between the two languages, take the

Fig. 6 Excitation of 2p2h states by the contact term (a) and the pionic current (b)

Fig. 7 An example of the excitation of a 2p2h state by the correlation current

specific example of the Pauli blocking effect which was derived in the virtual meson theory. We will restrict ourselves to one single contribution shown in Fig. 8. Can we trace back this effect in the cross-section? To recover the Pauli blocking effect in the cross-section we perform a cut in the graph of Fig. 8 so as to put the intermediate state on shell. By doing so we obtain a contribution to the photoproduction process in the quasi-elastic region, with 1p1h excitation. Near threshold there is indeed a strong influence of the Pauli blocking effect which should quench the nuclear photoproduction cross-section as compared to the free one [19].

For the quasi-deuteron effect we proceed in the opposite way, going from the cross-section to the meson exchange. The quasi-deuteron graph of Fig. 6b, for instance, is the imaginary part of a contribution to the Compton amplitude shown in Fig. 9 and its physical meaning becomes apparent. It is the same graph as that of Fig. 8, but now the exchanged pion has scattered on a neighbouring nucleon before getting reabsorbed by the first one. The effect of the quasi-deuteron cross-section can therefore be seen as a manifestation of the distortion of the pion cloud by the environment and we have seen that this has a large effect. Similarly we expect also a large influence on the r.m.s. charge radius. To get an idea of the change in the r.m.s. charge radius of the proton inside the nucleus we use the bremsstrahlung weighted sum rule (6). The difference from the corresponding relation for the free proton gives the r.m.s. radius increase. Above the pion threshold the free nucleon cross-section and the nuclear one per nucleon are close. Most of the r.m.s. radius modification then comes from the region below the pion threshold

$$\Delta\langle r^2 \rangle = \langle r_{\text{eff}}^2 \rangle - \langle r^2 \rangle_p \approx (3/e^2) \int_0^{m_\pi} d\omega \sigma_{\text{QD}}(\omega)/\omega \approx 0.6 \text{ fm}^2.$$

[expression (6) has been multiplied by four to account for the fact that the np distance in the quasi deuteron is twice the distance which enters in the electric dipole operator].

Fig. 8 A contribution to the e.m. polarizability of a nucleon inside a Fermi gas. The dash-dotted line shows the cut which puts the intermediate state on shell.

Fig. 9 A contribution to the Compton amplitude. The dash-dotted line shows the cut which leads to the photoabsorption process of Fig. 6b.

This represents a large increase, nearly a doubling of the free value $\langle r^2 \rangle_p = 0.65 \text{ fm}^2$. Is this value reasonable? Can we understand such a large effect? After all, a large increase is not surprising since the effective r.m.s. charge radius of the proton in the nucleus should feel the internucleon distance. If only 1/10 of the proton charge is transferred to a neighbouring neutron, at a distance of $\approx 1.8 \text{ fm}$, the r.m.s. radius increases by $\approx 0.3 \text{ fm}^2$.

Is this phenomenon the origin of the missing strength in the ee' experiments? We cannot be certain because the previous estimate is made in the static and long wavelength limit, which does not strictly apply to the actual experiments where the momenta range from 1 to 2.7 fm^{-1} and the frequencies from zero to $\approx 200 \text{ MeV}$. However, it is plausible that an effect of this origin exists and the magnitude should be qualitatively correct. It has nothing to do with a swelling of the quark core.

Strong experimental support in favour of this type of interpretation is provided by the following arguments. The effect that we have suggested concerns the *isovector* radius. The coupling of the photon to two pions is of an isovector nature; the isoscalar coupling proceeds instead through three-pion exchange. An increase in the isovector radius implies that the isovector form factor becomes, at finite momentum, quite different from its free value and hence from the isoscalar one which is not renormalized by this mechanism. Hence the coupling of the neutron to the photon, $G_s(q^2) - G_v(q^2)$, which is small for a free neutron, becomes appreciable in the nuclear medium. The photon '*sees*' the neutrons in a nuclear environment. This phenomenon is apparent in the experimental data, as the longitudinal response of ^{48}Ca is definitely different [7] from that of ^{40}Ca , which would not occur if G_s and G_v were equally renormalized. We have fitted [20] the experimental data, in the region of large momentum transfer, where the influence of collective effects should vanish, with the following inputs. The longitudinal cross-section is proportional to the quantity $Z(G_s + G_v)^2 R_p^0 + N(G_s - G_v)^2 R_n^0$, where R_p^0 and R_n^0 represent the responses of the free proton and neutron gases. The large renormalization of G_v in the nuclear medium provides a natural explanation of the difference between the ^{48}Ca response as compared to that of ^{40}Ca . In order to fit the experimental data, we have taken for the isoscalar form factor the free value $G_s(q^2) = (1 + q^2/\lambda_s^2)^{-2}$ with $\lambda_s^2 \approx 23 \text{ fm}^{-2}$, and for the isovector one a similar expression with an adjustable value of λ_v . A fit to the data is obtained with the value $\lambda_v^2 \approx 4 \text{ fm}^{-2}$ or equivalently an isovector radius $\langle r^2 \rangle_v = 3 \text{ fm}^2$. This increase is compatible with our previous estimated increase $\Delta\langle r^2 \rangle_v = 2\Delta\langle r^2 \rangle_p \approx 1 \text{ fm}^2$.

I consider the isotope effect to be an indication in favour of the change in r.m.s. radius associated with the two-pion exchange part of the coupling of the photon. It will be interesting to extend the exploration to the region of lower momenta, where the collective effects are important and the proton and neutron responses mix.

To summarize, I have proposed here a pionic interpretation of the swelling effect. The nucleon acquires an effective r.m.s. charge radius much larger than the free value, from the spreading of the charge to the neighbouring nucleons or to the pion cloud.

This type of swelling affects exclusively the isovector swelling. It does not concern the quark core, which remains unchanged. This picture has two merits. Firstly it explains the differences observed in the longitudinal responses of the two isotopes ^{40}Ca and ^{48}Ca . Secondly, a swelling of this nature, which does not touch the quark core, explains why the excitations of this core do not seem to be changed in the nuclear environment. For instance, the Δ resonance is found at the same excitation energy in the nucleus and for a free proton. A large swelling of the core instead should normally lower this energy [21] and one has to invoke compensating effects to justify this fact. Of course our interpretation does not preclude the existence of a genuine (but more modest) modification of the core radius. But it makes unlikely its observation in the ee' longitudinal response, which is more sensitive to pionic effects.

The question that I want to raise now concerns the rescaling assumption in the EMC effect. Does the pionic swelling, which represents a change in scale, entail the rescaling? I believe it does not, as rescaling implies a change in scale for the valence quarks, which does not take place here.

So in order to have rescaling as the cause of the EMC effect we have to accept the idea that this effect is totally disconnected from the missing strength problem in ee' scattering. Personally I believe instead that the two effects are two facets of the same physical reality, that of pions in nuclei. The second choice is esthetically more appealing and likely to be closer to reality.

REFERENCES

- (1) R. Jaffe, F. Close, R. Roberts and G. Ross: Phys. Lett. **134B**, 449 (1984).
- (2) C.A. Garcia Canal, E.M. Santangelo and H. Vucetich: Phys. Rev. Lett. **53**, 1430 (1984).
- (3) S.V. Akulinichev, S. Shlomo, S.A. Kulagin and G.M. Vagradov: Phys. Rev. Lett. **55**, 2239 (1985).
- (4) C. Llewellyn Smith: Phys. Lett. **128B**, 107 (1983).
M. Ericson and A.W. Thomas: Phys. Lett. **128B**, 112 (1983).
- (5) V.R. Pandharipande: Nucl. Phys. **A446**, 189c (1985).
- (6) M. Oka and R.D. Amado: University of Pennsylvania preprint UPR-0295T (1986).
- (7) P. Barreau et al.: Nucl. Phys. **A402**, 515 (1983).
Z. Meziani et al.: Phys. Rev. Lett. **52**, 2130 (1984) and **54**, 1233 (1985).
- (8) J. Noble: Phys. Rev. Lett. **46**, 412 (1981).
- (9) C.M. Shakin: Nucl. Phys. **A446**, 323c (1985).
- (10) M. Ericson and M. Rosa-Clot: preprint CERN-TH 4420/86 (1986).
- (11) M. Rosa-Clot and M. Ericson: Z. Phys. **A320**, 675 (1985).
M. Ericson: Lectures at the School on Nuclear Dynamics, Dronten, 1985, preprint Lycen/8577 (1985).
- (12) G. Höhler and E. Pietarinen: Phys. Lett. **53B**, 471 (1975).
- (13) P. Weiner and W. Weise: Phys. Lett. **159B**, 59 (1985).
- (14) E. Oset, R. Tegen and W. Weise: Nucl. Phys. **A426**, 456 (1984).
- (15) S. Theberge, G. Miller and A.W. Thomas: Canad. J. Phys. **60**, 59 (1982).
- (16) T. Ericson and J. Hüfner: Nucl. Phys. **B47**, 405 (1972).
- (17) J. Delorme, M. Ericson, A. Figureau and C. Thévénat: Ann. Phys. **102**, 273 (1976).
M. Ericson: Progr. Part. Nucl. Phys. **1**, 67 (1978).
- (18) J. Ahrends: Nucl. Phys. **A446**, 229c (1985).
- (19) W.M. MacDonald, E.T. Dressler and J.S. O'Connell: Phys. Rev. C **10**, 455 (1979).
- (20) W.M. Alberico, P. Czerski, M. Ericson and A. Molinari: in preparation.
- (21) J. Noble: Abstract A17, p. 82, 11th Europhysics Divisional Conf. Paris, July 1985.