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T.E.O. Ericson

CERN - Geneva

A B S T R A C T

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T.E.O. Ericson

Theoretical Physics Division, CERN, 1211 Geneva 23, Switzerland

The present evidence for pion contributions to magnetic moments, transitions and form factors is selectively reviewed with emphasis on the crucial evidence for specific effects. The theoretical basis will be discussed with special emphasis on the physical picture which emerges in the long wavelength limit.

Professor Arima /1/ and Professor Stone /2/ have already provided the general setting within which we should discuss magnetic effects in nuclei beyond the nucleon description. Pions have a special significance for magnetic nuclear physics for two reasons /3/.

1. The gyromagnetic ratio of the pion is much larger than the nucleon one:

$$\frac{e}{2m_{\pi}} \approx 7 \frac{e}{2M} \quad (1)$$

As a consequence small pion admixtures to a nucleonic wave function give anomalously large contributions.

2. The pion is the main nuclear constituent beyond nucleons and plays a major rôle in nuclear binding: for that reason alone one expects its rôle to be of special importance. The magnetic effects associated with pions have mainly isovector character, since the pion has isospin $I = 1$. But will not a pion description conflict with, say, a quark-gluon description? Is it not arbitrary to insist on the pion aspects? Not at all. In different parts of nuclear physics, we must pick the relevant degrees of freedom for the most efficient and economical description of the physics on the scales of the problem. It should be a sobering thought that nuclear spectroscopy to an important extent never needs nucleons: effective bosons, quasi-particles and collective degrees of freedom provide a powerful framework with excellent physics. At high energies with short wavelengths, quarks and gluons are appropriate. In nuclear pion physics, the "quark-clusters" which describe the physics very well are the π 's, N 's, Δ 's and some ρ 's. These degrees of freedom summarize an impressive amount of what we might call "non-perturbative QCD" in nuclei. The quarks are there, but they hide inside these building blocks.

It is appropriate to ask how well nucleons provide an approximate description of nuclei. For this we have to eliminate the model dependence on nuclear wave functions, which are a curse on precision tests of this kind.

It has been suggested by Karl et al. /4/ that nucleon magnetic moments in nuclei are modified substantially by quark deconfinement. A comparison /5/ of the allowed Gamow-Teller $\beta^\pm$ transitions to the corresponding M1 transition from the supermultiplet 1^+ states in the $A = 12$ system tests this to a few per cent accuracy: there is a near identity between the β -decay and magnetic moment spin matrix elements! In the impulse approximation for nucleons only, the operator is:

$$\beta^\pm : g_A \sum \vec{\sigma}_i \cdot \vec{\tau}_i^\pm \quad (2)$$

$$M1 : \frac{g_p - g_n}{2} \sum \sigma_{iz} \tau_{iz} + \text{small orbital corrections} \quad (3)$$

After a 5% correction for well-defined and well-understood pion effects the agreement is excellent. Thus nucleons are a pretty good start on the few per cent level.

From the viewpoint of particle physics the pion is a very special particle with a nearly vanishing mass (certainly not small for a nuclear physicist!). This is associated with the so-called chiral limit of quantum chromodynamics (QCD) in which u and d quark masses are very small. For $m_{u,d} \rightarrow 0$ it follows that the pion mass $m_\pi \rightarrow 0$ in the same limit. This is referred to as the soft limit of pions: as $m_\pi \rightarrow 0$ the Compton wavelength $1/m_\pi$ becomes large and the details of hadron structure cannot be explored. In this soft limit, somewhat analogous to the Thomson limit for soft photons, a number of theorems apply. In particular, the threshold pion scattering amplitude becomes universal from any target and pion photoproduction also becomes universal from any target. In practice these theorems apply nicely to elementary nucleons, which are nearly pointlike on the scale of the pion Compton wavelength. For nuclei there are gross violations. The explanation is that scales are not preserved in the soft limit for nuclei. The pion Compton length is small on the nuclear scale, while the pion mass is large compared to typical nuclear excitations. If the limit is taken in such a way that scales are respected, soft theorems hold in nuclei, but at the level of individual nucleons.

The second major chiral consequence is that that pion physics is largely equivalent to axial currents, i.e., to spin-isospin physics. This is of course particularly relevant to magnetic moments and transitions, since these concern axial operators.

We all know that electromagnetic currents are conserved: $\partial_\mu J^\mu = 0$. The axial current A^μ is nearly conserved; if it were strictly conserved, however, the pion would never decay. The violation follows from chiral symmetry and is due to the pion field ϕ . This is referred to as the partially conserved axial current (PCAC):

$$\partial_\mu A^\mu = -f_\pi m_\pi^2 \phi \quad (4)$$

where f_π is the pion decay constant. From PCAC and with a natural locality assumption on the level of individual nucleons (axial locality) one finds that the space and time parts of the axial current are related to p- and s-wave pions respectively:

$$\vec{A}^+ \approx \sum \vec{\sigma}_1 \tau_1^+ + \text{exchange} + \Delta \quad (\text{p-wave pions}) \quad (5)$$

$$A_0^+ \approx \sum (\text{velocity terms}) + \text{exchange} \quad (\text{s-wave pions}) \quad (6)$$

The axial current operator acts exactly like an external s- or p-wave pion on the nucleon level: axial currents and pions are equivalent!

Let us now look at magnetic operators from this viewpoint. First, the pion is an isovector, so that we only make statements about the isovector magnetic operators. Furthermore, we make at the moment no statements concerning orbital or convection currents, even for the pion. These will be discussed later. On the level of the ΔN contributions to the magnetic moment operator we have the following scaling relation including exchange currents:

$$\vec{m}_V(\vec{x}) = \frac{e}{g_{\pi NN}} \frac{\mu_V}{f_\pi} \vec{A}_3(\vec{x}) \quad (7)$$

relating the isovector axial and magnetic moment operators. On the nucleon level we easily see the power of this relation. The pion s- and p-wave photoproduction on the nucleon are E1 and M1 transitions (Fig. 1). The E1 production of s-wave pions follows from simple classical considerations or from soft pion arguments: it involves essentially the replacement of $\vec{\sigma} \cdot \vec{V}$ by $\pm e \vec{\sigma} \cdot \vec{\epsilon}$ by minimal electromagnetic coupling. The M1 excitation of the Δ can be looked at as a quark magnetic spin-flip transition. The parameter-free scaling relates the M1 amplitude $M_{\gamma\pi}$ to the πN scattering amplitude $f_{\pi N}$ for p-wave pions:

$$M_{\gamma\pi} = e \frac{\mu_V}{g} \frac{\omega}{q} f_{\pi N} \quad (8)$$

As is seen in Fig. 2, this relation is very accurate.

Fig. 1. The dominant E1 and M1 contributions to the nucleon photo cross-section.

Fig. 2. The dominant M1 amplitude scaled from πN scattering according to Eq. (3) (solid line) compared to its empirical value (dots).

The best nuclear tests of axial locality and chiral symmetry involve β -decay $0^+ \rightarrow 0^-$ transitions, which are changed by large factors due to s-wave pion rescattering. However, this falls outside the topic of my talk.

The intrinsic axial contributions of the axial current to magnetic operators in nuclei are conceptually important but are not the dominant ones apart from the

impulse approximation terms from the individual nucleons. They appear implicitly from transition and convection currents associated with pion exchange as I will now discuss. Let me first recall some of the crucial features of the one-pion exchange potential (OPEP) [7]. This potential is closely similar to the magnetostatic potential associated with the interaction of two bar magnets of magnetic moment $\vec{\mu}_1$ and $\vec{\mu}_2$. In this case the photon exchange leads to a hyperfine interaction (as appropriate for this meeting) with a central Fermi contact interaction and a tensor potential acting on states with $\ell = 0$. As seen in Fig. 3 the pion exchange closely parallels this situation. The pion exchange couples to each of the nucleons by an axial dipole strength $f/m_\pi \vec{\sigma}_{1,2}$, instead of $\vec{\mu}_{1,2}$. The coupling is longitudinal instead of transverse in the gradient operator; further the pion mass gives rise to a Yukawa field $\exp(-m_\pi r_{12})/r_{12}$ instead of $1/r_{12}$. As a consequence of this axial dipole-dipole coupling there is thus a (weak) central potential corresponding to the Fermi contact interaction and a very strong tensor potential. It is this last feature that governs nuclear physics to a considerable degree.

Fig. 3. The analogy between the magnetic dipole-dipole interaction and the pion axial dipole-dipole interaction.

How quantitative is the OPEP description in a situation that can be controllably tested? The deuteron is a system which we can explicitly calculate using the OPEP potential, and for which we know both wave functions and physically measurable quantities to very high precision, typically per cent to per mille precision. If we insist on the experimental binding (essential, since small variations in the fine-tuning of this quantity will falsify all scales in the problem) we can integrate the coupled S- and D-state equations from infinity inwards; by a small trick one can obtain "wave functions" regular in the origin (Fig. 4). There is no free parameter in the problem. The S-wave function $\tilde{u}$ is nearly identical to the empirically derived one, and to the most sophisticated approaches on the market. The D-wave function is asymptotically excellent but

Fig. 4. The deuteron S and D wave functions generated from iterated OPEP /8/ and compared to those of the Paris potential as well as empirical deduced values.

somewhat too strong at about 1.5 fm reflecting the too strong OPEP tensor force (which in a less primitive approximation is tempered by ρ -exchange, for which the leading contributions are also parameter-free!). Even with this simple approach one finds that a number of basic parameters are reproduced on a level of five per cent or better: the effective range, the asymptotic D/S ratio, the quadrupole moment and radius. It is easy to improve this down to the few per cent level model-independently; there one reaches experimental uncertainties from particle physics due to the πNN coupling constant and its form factor.

As an example of the quantitative understanding of these quantities, consider the interaction density for the quadrupole moment (the interaction density for the asymptotic D/S ratio is very similar). This quantity is approximately proportional to $r^3 V_{\text{tensor}}(r) \tilde{u}(r)$ and is given in Fig. 5. There are very small contributions inside 1 fm, where quarks may play a rôle; the -6% deviation between the OPEP value and that for the Paris potential is quantitatively and model-independently associated with ρ exchange.

In view of this understanding, it can be said that the deuteron is the most quantitative example of nuclear pion physics.

Fig. 5. The deuteron quadrupole interaction density for OPEP compared to the same quantity for the Paris potential /8/.

These results for the deuteron are echoed in other forms throughout the periodic system although with decreasing precision as nuclei get heavier. In particular, the intrinsic deformation of the α -particle (i.e., the coupling of the four nucleon spins to $S = 2$ with an orbital $L = 2$ component giving $J = 0$), the nuclear binding as well as the tensor correlations in heavy nuclei are all of pionic origin dominated by the tensor potential.

Let us now turn to the main magnetic effects associated with pions. They are the following /9/.

1. The effects arising from the exchange of a charged pion between a neutron and a proton. This is the counterpart of the E1 pion photoproduction on a nucleon to which should be added the current from the pion in flight (see Fig. 6). This contribution is obtained by minimal electromagnetic coupling to the $\pi N \vec{\sigma} \cdot \vec{\nabla}$ interaction. The main magnetic effects of these nuclear contributions are two:

- the Sachs moment which is the effect of the convection current of pion exchange; it renormalizes the isovector orbital g-factor δg_λ , so that it has opposite sign for neutron and proton;
- an intrinsic two-body effect yielding a transition operator between np singlet and triplet states with $\Delta I = 1$:

$$^1S_0 \leftrightarrow ^3S_1 + ^3D_1$$

Fig. 6. The one-pion exchange current.

2. The Δ contribution to the exchange current, which in particular renormalizes the nucleon axial coupling constant g_A and the spin g -factor g_s .

The clearest evidence for the contributions of the two-body ${}^3S_1 + {}^3D_1 \leftrightarrow {}^1S_0$ comes from the electrodisintegration $ed \rightarrow e'(np)$ at threshold which is a spin-flip M1 transition, which has been explored in great detail in the last few years and where the effects are huge. The closely related M1 radiative capture process $np \rightarrow \gamma d$ for thermal neutrons has played a very important historical rôle, since Brown and Riska /10/ found the explanation of the 10% model-independent discrepancy with nucleons - only in terms of exchange currents due to the pion. Figure 7 shows how modern experiments permit the exploration of the distribution of this transition magnetization up to quite large momentum transfer allowing a spatial resolution of the order of $1/q$ of less than 0.5 fm. It is apparent from the figure that a description in terms of nucleons - only fails very badly. In view of the high degree of resolution one may wonder whether this is a quark or a pion effect. Until recently it was usual to calculate the meson exchange contribution assuming simply one-pion exchange or simple generalizations of it. This procedure, although it gives reasonable results, was always criticized, since at short distances one can have contributions from, say, 5π or 12π or 137π exchange in an uncontrollable way, or quark exchange.

However, we nowadays know the NN interaction so well that we can empirically deduce the mass distribution for the exchange of π or ρ quantum numbers from the NN amplitudes. For these masses one can then make minimal couplings, as if these were pointlike and elementary. The curve which closely follows the experimental points in Fig. 7 has been obtained without parameters in this way. It is furthermore easy to construct the corresponding interaction density (Fig. 8) for different momentum transfers. One sees clearly that the region inside 1 fm is unimportant. The explicit lesson from this approach is:

- a long-range interaction density with short-range suppression;
- a very dominant pion contribution;
- an empirically small " ρ " contribution.

Fig. 7. Meson exchange effects in backward deuteron electrodisintegration near threshold as a function of momentum transfer /11/.

Fig. 8. The OPE transition density in the process $ed \rightarrow e'np$ at threshold /12/.

The same magnetic meson-exchange contribution coupling neutron-proton $^3S_1 + ^3D_1 \rightarrow ^1S_0$ states appears also dominantly in the $A = 3$ magnetic form factors as seen in Fig. 9. This is also by far the most important meson-exchange contribution in the M1 isovector transition in ^{12}C . We can therefore conclude that this feature of the magnetic effects in nuclei from pions is very well established.

Fig. 9. Meson exchange current effect in the $A = 3$ system as a function of momentum transfer /13/.

The remaining effects due to pions are much more qualitative. Inside an uncorrelated Fermi gas picture the orbital g -factor is modified by the central exchange potential. The physics of this effect is essentially that the current $\vec{J}_p = e\vec{v}_p$ of a proton has to be supplemented by the pion current $\vec{J}_\pi = e\vec{v}_\pi$, in which the pion "jumps ahead" to a neutron carrying the proton momentum. The pion is then quite effective in the magnetic moment, since its mass is seven times less than that of the nucleon. In spite of its intrinsic weakness the central OPEP gives the most important contribution with $\delta g_\lambda^\pi(p) \approx 0.15$ and $\delta g_\lambda^\pi(n) \approx -0.1$, in rough agreement with experiments as discussed by Prof. Stone /2/. The uncertainty comes from the fact that there are large and cancelling contributions, both from tensor correlations as Prof. Arima /1/ has discussed, and from the corresponding effect for ρ -mesons. There is little doubt about the importance of the pionic terms; one can certainly improve the description beyond this level in this conceptually very interesting quantity, but I doubt very much that any reliable statement concerning quark physics can be extracted from δg_λ .

The evidence for the Δ -effects in magnetic moments and transitions is equally qualitative. In this case the connection to pions goes via the axial current and "axial locality": the contributions and operators are expected to be identical to those for axial transitions such as Gamow-Teller β -decay matrix elements, and the space part of the axial current is associated with p-wave pions. The explicit exchange contributions associated with the rescattering of the pion on one nucleon before its absorption on a second one stand out more clearly in β -decay than in magnetic moments, where they are overshadowed by the previously discussed $^3S_1 + ^3D_1 \leftrightarrow ^1S_0$ effect. However, both for the magnetic moment in $A = 3$ and for the $np \rightarrow \gamma d$ M1 transition in the deuteron, agreement with experiments is improved when the pion-exchange current contribution is included /6/. For magnetic moments it stands out most clearly as a renormalization of the spin gyromagnetic ratio g_s in heavy nuclei in analogy with the effective axial coupling constant. The physics of this renormalization is quite interesting and is due to the predominant dipole interaction of the pion in the nuclear medium /14/.

The p-wave πN scattering in the nuclear medium leads to a susceptibility χ for the propagation of a pion wave. This is just like electromagnetic dipole scattering. In both cases there will be induced dipoles in the medium. An ordinary dipole in a dielectric will have its strength changed by the induced dipoles in the little cavity in which it is sitting; this is the consequence of the local effective field. In precisely the same way the axial dipole with strength g_A of a nucleon is renormalized:

$$g_A^{\text{eff}} = \frac{g_A}{1+g'\chi} \quad (9)$$

In real life, the susceptibility has two main contributions: χ_N from the repeated $\vec{\sigma} \cdot \vec{\nabla}$ -coupling to the nucleons and χ_Δ from the virtual excitation of the Δ . The classical value for this parameter is $g' = 1/3$, the so-called Lorentz-Lorenz effect. This value is observed in low energy πN scattering of physical pions on nuclei, which is Δ dominated. The value of this parameter for nucleons in a situation of near-vanishing pion energy ω is $g' \approx 0.7$ to 0.8 : if this were not so, the system would be too attractive and pion condensation would result. As a consequence, $g'\chi$ should be replaced by $g'_N \chi_N + g'_\Delta \chi_\Delta$. Both the nucleon contributions and the Δ contributions are important. As Prof. Arima has emphasized, tensor correlations are the main contributions; this is another way of talking about the polarization of the nuclear medium. In both cases the renormalization occurs. This is exactly the conclusion that has also been reached in discussions of pion condensation. This phenomenon has little to do with "pions" as physical particles. It corresponds to a spin-isospin ordering in layers of nuclear matter, which is provoked mainly by the tensor potential. The corresponding spin-isospin excitations have the quantum numbers of the pion (see

Fig. 10). Typically, χ_N is $(2-3)/\chi_A$. As a consequence of this renormalization, the spin gyromagnetic factor g_s is systematically quenched, like g_A is quenched, in heavy nuclei (Fig. 11). The magnitude of this effect is well consistent with pion physics. The importance of such renormalizations in nuclei goes well beyond magnetic moments.

Fig. 10. The spin-isospin alternating layer structure.

Fig. 11. The effective spin g-factor determined from M1 transitions measured in (e,e') reactions [15].

In conclusion, pion physics appears in a wide variety of nuclear phenomena, usually associated with pion-like quantum numbers and axial matrix elements. In magnetic phenomena the operators associated with the direct coupling to nucleons are well established from light elements and, the effects on g_l and g_s are qualitatively correctly described by pion physics. Alternative explanations for such effects must be quantitatively consistent with well-established aspects of nuclear pion physics.

References

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