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Recent Developments in Nuclear Physics*

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Abstract

Nuclei exhibit features that are described in superficially contradictory terms according to the different degrees of freedom that are excited by probes of different scale in space and in time. After giving some examples I concentrate on the hadron degrees of freedom such as the nucleon, the pion and the Δ isobar. These are the effective degrees of freedom on the level of intermediate resolution: about 0·5–1 fm in distance and correspondingly in time. A prime example is the deuteron which has a nearly model-independent description in terms of pion physics to very high precision. In nuclear matter the pion propagates in close analogy to the propagation of light in a dielectric. This permits the explanation of a number of features in nuclei related to the chiral symmetry limit in which the pion mass vanishes. A consequence of this description is the analogy of the equations for the pion and its effective field with the Maxwell equations for a dielectric. A pionic collective mode should appear strongly and with characteristic properties for a well chosen probe. It is difficult to explore its properties directly and in particular physical pions are not useful for this purpose. I will discuss different alternatives involving 'virtual pion beams'. There is recent evidence for such a collective state in forward charge exchange reactions throughout the periodic system.

1. Scales in Nuclear Physics

One of the most remarkable and fascinating aspects of nuclear physics is that we meet a series of descriptions of nuclei which at first glance appear totally contradictory, but which only describe different approximation schemes to the same many-body problem. The situation is a consequence of the fact that we use probes with very different resolutions in space and in time for the exploration of the nucleus. With each level of resolution different effective degrees of freedom are appropriate for the quantitative description of the nuclear phenomena.

A brief examination of some of the present frontier areas in nuclear physics brings out this situation very well. First, probes with low resolution explore nuclei either at wavelengths large compared with the nuclear radius, or for times long enough to permit the study of closely packed and sharp nuclear states. This situation emphasises the behaviour of the nucleus as a semi-classical system. A good example is the spectacular observation (Twin *et al.* 1986) of the rotational spectra of superdeformed nuclei with angular momentum of about $I = 60$ (Fig. 1). The description of this situation hardly requires that nuclei consist of neutrons

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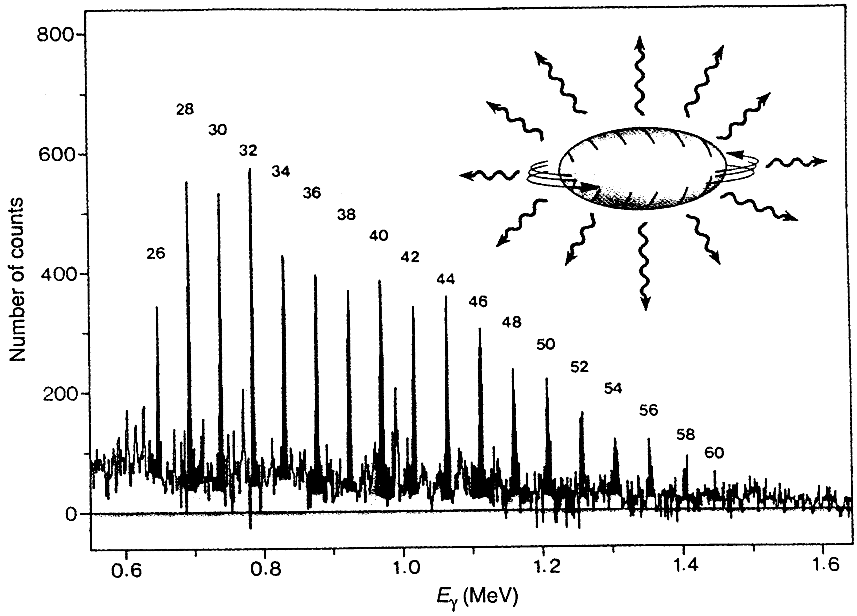

Fig. 1. The superdeformed rotational spectrum of ^{152}Dy (Twin 1991).

and protons: it is quite enough, and very interesting at that, to analyse such a structure in terms of the quasi-classical properties associated with the stability of a rotating liquid drop.

While this may be satisfying in this limit, we are well aware that valence nucleons moving rather independently in a smooth potential well with a shape roughly corresponding to the nuclear density distribution provide an excellent description of a large number of nuclear properties and nuclear spectra at long wavelengths. In this case, however, we are just scratching the nuclear surface and it has been a long standing issue whether in the deep nuclear interior of high density such a description has much meaning, in spite of its success. It is therefore exciting that electron scattering presently probes the nuclear charge distribution to such a precision that we can directly measure the wavefunction of a single valence proton. Fig. 2 gives the empirical wavefunction in r-space for the 3s proton state in Pb: one sees clearly the three nodes which are its characteristic signature. The remarkable thing is that we also see a sharp peak, as expected from the potential description, inside a region of only 1 fm, far smaller than the Pb radius of 8 fm and in a region of maximal matter density. Such experiments tell us that the shell model has a very literal meaning and that its degrees of freedom work very well up to moderate resolution even at high nuclear density.

As we move towards a higher resolution of the nucleus, the nucleon degree of freedom comes more strongly into focus, but so do other hadrons, in particular the pion as the main carrier of the nuclear force and the Δ isobar as the first excited state of the nucleon. This is the regime of hadron physics and it will be the main topic of this paper. We meet much beautiful physics here and some of the most quantitative predictions of nuclear physics belong to this field.

Amazingly, we now know there is hardly any trace of quark physics at the hadron level, although we know from particle physics that quarks are the ultimate constituents of nucleons and pions and thus of the nucleus. Because of that there is a strong effort under way to explicitly demonstrate the importance of quarks also in nuclei.

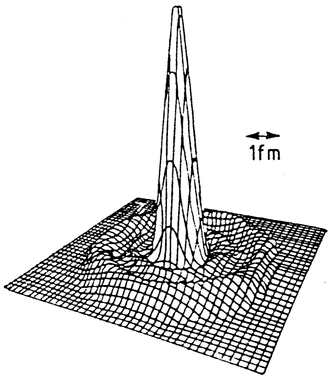

Fig. 2. The observed 3s wavefunction in ^{208}Pb (Frois and Laget 1983).

The first of these efforts goes under the name of the EMC effect and related phenomena (see e.g. Thomas *et al.* 1993, present issue p. 3). In essence, the unexpected finding is that when nucleons in a nuclear environment are explored with high resolution using deep inelastic muon scattering, they appear to grow bigger with their quarks deconfined. There has been an important debate about how we should view this phenomenon most profitably, in particular whether we should look at the deconfined quarks as part of a modified nucleon-pion field or as associated with binding effects. What is clear, however, is that the nucleons, although modified, maintain the essence of their identity in this regime. The EMC effect therefore corresponds to quark perturbations of the hadron regime.

The way to study quarks in nuclei is obviously to make the quark aspect totally dominant. This is exactly what is now being attempted at the SPS at CERN and at the AGS at Brookhaven by going to yet another scale. If the distance between the nucleons is compressed to less than 0·5 fm and the nucleus is simultaneously heated to a temperature of about 200 MeV, rather general arguments indicate that the quarks should be liberated into a new phase of matter, a quark–gluon plasma, a state that is likely to have been typical of the early universe following the big bang. The experimental difficulty is on the one hand to achieve such densities and temperatures during a sufficient time for thermal equilibrium, and on the other hand to obtain an unambiguous signal for the formation of such a plasma. The road to achieve the conditions is by colliding extended matter, i.e. nuclei, with each other at relativistic energies. The collision should be head-on for maximal conversion of energy into heat. Such experiments are being carried out very systematically. At present it appears that the necessary energy density and temperature have been reached, or nearly so. The problem is the signal. Candidates are anomalously high production of charmed particles (J/ψ) or strange antibaryons like anti-Λ or anti-Ξ, which are the consequence of quark deconfinement. The present excitement stems from the fact that there is indeed an apparent signal, as shown in Fig. 3. The key issue is now to demonstrate that this is not an unwanted background.

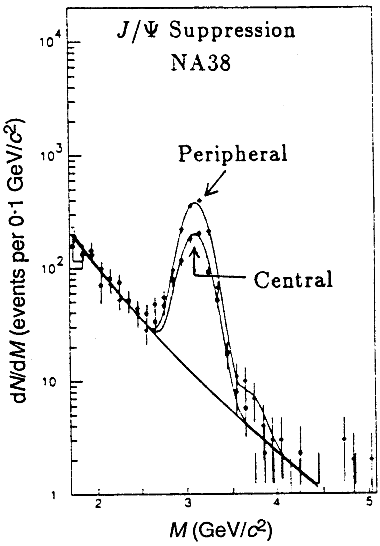

Fig. 3. The suppression of the J/ψ signal, which may indicate quark–gluon plasma formation (Leitch *et al.* 1991).

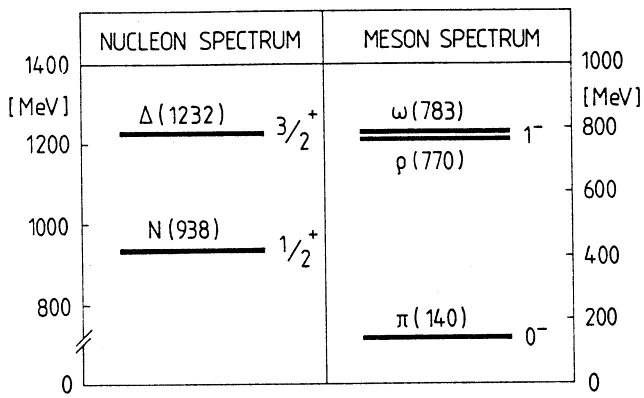

Fig. 4. The principal low energy hadrons with their mass in MeV.

In the present paper I will explore only one of these areas in more depth, namely the one of nuclear pion physics. In this limit the main nuclear constituents are the nucleons, the pion and the Δ isobar. Fig. 4 gives their schematic properties and relative energies.

2. Role of Pions in Nuclei

The pion has a very special role in nuclear physics. Not only is it the lightest by far of all the mesons, but it is the main agent of nuclear binding. This alone sets it apart from other nuclear probes. In addition, however, its low mass is deeply related to a special symmetry called chiral symmetry, which in practice corresponds to a limiting situation in which the pion mass vanishes. The zero frequency collective mode of this symmetry, its Goldstone boson, is the pion. The chiral symmetry introduces powerful constraints, even theorems, concerning the interaction of the pion with nucleons and with nuclei. The pion is thus a full partner in the nuclear many-body problem: this sets it deeply apart from other probes like electrons and kaons. Indeed, pion physics permeates the whole of nuclear physics on the microscopic level. Examples are

Nuclear interactions:

— the key part of the nucleon–nucleon interaction, the tensor force,
— the quantitative properties of the deuteron,
— nuclear tensor correlations,
— the bulk of 3-body forces,
— the newly observed pionic collective mode in nuclear matter.

Axial currents, which are identical to pion physics within certain general assumptions

— spin–isospin excitations,
— β-decay and μ-capture as well as nuclear neutrino physics,
— magnetic moments and transitions, all including exchange currents.

Physical pions:

— Δ isobar dynamics,
— pair correlations,
— physics of the pion as a probe of the nucleus,
— tests of isospin purity of nuclear states.

Nuclear chiral symmetry:

a heavy constraint and guiding principle in nuclei.

Clearly I cannot survey all these varied aspects. I will thus concentrate on four topics that have special actuality and importance:

1. Pion dominance of the deuteron;
2. Pions, the nuclear Lorentz–Lorenz effect and the pion effective field;
3. Axial currents and nuclear chiral symmetry;
4. Evidence for a new collective pion mode using 'virtual pion beams'.

3. Pion Dominance of the Deuteron

There is currently much discussion about the modifications of the nucleons and the nuclear effective degrees of freedom in terms of quarks. It is then important

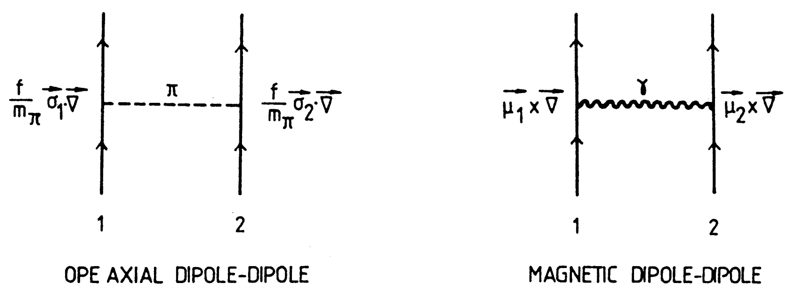

Fig. 5. The analogy between the magnetic dipole–dipole interaction and the pion axial dipole–dipole interaction.

to establish quantitative limits on the tolerable extent of such modifications within a pionic picture. At the same time this gives an occasion for a much more profound insight into the hadronic approach to nuclear physics. The central ingredient in the long-range nuclear interaction is the one-pion-exchange (OPE) potential. It is remarkable that this potential has a close similarity to the hyperfine potential well known from the interaction between two magnetic dipoles in atomic physics (see Fig. 5).

In the magnetic case we have two axial dipoles interacting through a *transverse* dipole coupling of the two magnetic moments $\mu_{1,2}$ and with the exchange of a massless photon. This leads to the usual magnetostatic potential. In the pion case we also have two interacting axial dipoles $f/m_\pi \sigma_{1,2}\tau^\alpha$ but now interacting through a *longitudinal* coupling and with the exchange of a pion of mass m_π. This is why the two potentials have a nearly identical structure in the static limit:

$$V_{\mathrm{magn}} = (\mu_1 \times \nabla)(\mu_2 \times \nabla)\,(1/4\pi r)\,,$$

$$V_\pi = (f/m_\pi \tau_1^\alpha \sigma_1 .\nabla)(f/m_\pi \tau_2^\alpha \sigma_2 .\nabla)\,\exp(-m_\pi r)/4\pi r\,.$$

The πNN coupling thus occurs via nucleon *axial dipole moments* $m_{\mathrm{i}}{}^\alpha = f/m_\pi \sigma_{\mathrm{i}} \tau^\alpha$ in terms of the nucleon spin σ and isospin τ^α, in analogy with the magnetic moment μ_i. The pionic potential is damped by the usual Yukawa factor, but it has otherwise the same shape as the magnetic interaction. The key feature in the pionic potential is the strong tensor force; for the case of the magnetic hyperfine interaction in an atom, on the contrary, the tensor force is no more than a perturbation of the hydrogenic spectrum. The importance of this tensor force to nuclear binding comes out clearly in the case of the deuteron, which I consider to be the prime example of nuclear pion physics. It gives to high precision the quantitative

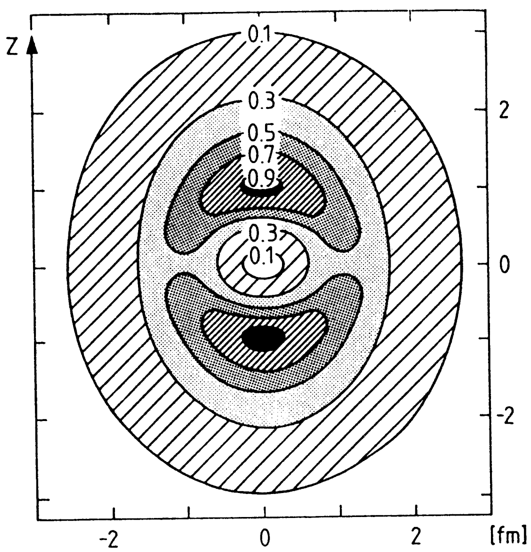

Fig. 6. Deuteron density contours derived from Paris wavefunctions (Ericson and Weise 1988).

details of the deuteron without free parameters (Ericson and Rosa-Clot 1985). Let me first show in Fig. 6 a density contour plot of the deuteron derived from standard deuteron wavefunctions: it shows that there is a very clear separation of the density, which is on the average concentrated around two centres. Clearly the deuteron has much similarity to a diatomic molecule.

In principle one might think that a quantitative description of the deuteron in terms of the OPE potential could be obtained from direct evaluation of the corresponding deuteron Schrödinger equation. This, however, completely obscures the issue for the simple reason that the deuteron is very nearly unbound; the scales in the problem and, in particular, the deuteron size are largely determined by the rather accidental, but crucial, value of the binding energy which results from a cancellation of two large numbers. It is essential to fine-tune the short-range interaction so as to impose the correct binding energy, i.e. to give correct asymptotic behaviour of the wavefunctions, while preserving regularity at the origin. Fig. 7 gives the result compared with one of the best wavefunctions on the market (Lacombe *et al.* 1980), as well as with the wavefunctions empirically deduced from electron scattering form factors (Locher and Svarc 1984).

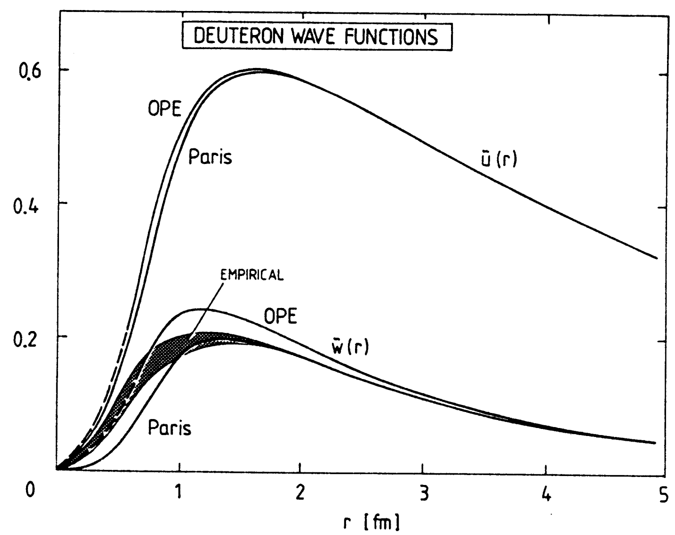

Fig. 7. The deuteron s and d wavefunctions from iterated OPEP, from the Paris potential and from those determined empirically (Ericson and Rosa-Clot 1985).

One notes at once that the spherical component of the deuteron described by the s wave is in nearly perfect agreement (to within a few %) with the results of the Paris potential. This means that it is in equally good agreement with the wavefunction directly deduced from electron scattering, which we have not plotted in the figure because it agrees so closely with our result that it becomes almost indistinguishable. All this is a consequence of the *iterated* tensor interaction, an interaction that does not contribute at all to the s state in leading order! Also the d wave comes out surprisingly well, although it is somewhat too strong in the inner region. This is so because the tensor potential for pure OPE is unrealistically strong at intermediate and short range. In real life it is tempered by 2π exchange and form factors, which can be included quantitatively. Clearly this zeroth-order approach is amazingly good and rivals the sophisticated Paris potential approach. Table 1 shows how accurately pion physics describes the deuteron properties in this simplest approximation:

Table 1. Key deuteron properties in the simplest pion approach

	Pion only	Experiment
Binding energy	Fixed by fiat	2·22457 MeV
Effective range	1·74 fm	1·765(5) fm
Asymptotic d/s ratio η	0·0276	0·0271(4)
Quadrupole moment	0·294 fm^2	0·2859(3) fm^2
Radius	1·94 fm	1·963(4) fm

One thus has a parameter-free description to within a few % which can easily be improved. This precision is so high that the experimental value of the πNN coupling constant is a practical limit for the description. This result is a heavy constraint on any claim for exotic features in the ground state of the deuteron, such as quark degrees of freedom, since there is no need for them.

All this success is mainly a consequence of the tensor force and its iterations at distances greater than 1 fm. The properties inside this region are essential for getting the binding energy correct, but they are not very essential for the rest of the deuteron properties. On the other hand, between 1 and 2 fm the OPE tensor potential is systematically weakened by the theoretically well understood 2π exchange with ρ meson quantum numbers, as well as by the fact that the πNN source is spread out in space; such contributions can be calculated model-independently by dispersive methods outside of about 0·6 to 0·7 fm, as seen in Fig. 8.

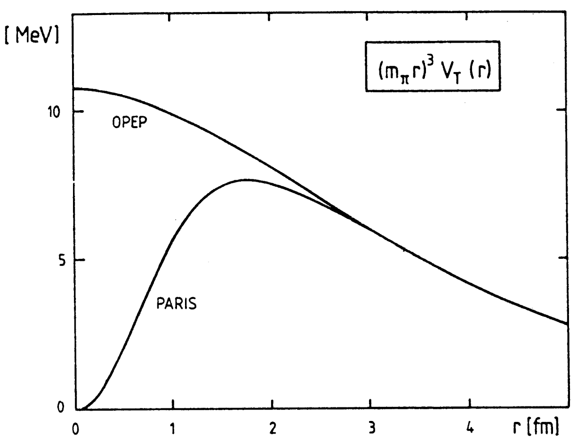

Fig. 8. The isovector tensor potential from OPE and from the Paris potential (Ericson and Weise 1988).

One can ask whether this insight also applies to electromagnetic and dynamic phenomena. Here a special feature of the pion is its large gyromagnetic ratio of about $7e/M$ nuclear magnetons. Consequently we expect the pion to deeply influence nuclear magnetic phenomena, since even a small component of the pion in the wavefunction will have a disproportionately large effect. This appears clearly in the threshold electro-disintegration of the deuteron, which has been carefully investigated at Saclay up to large momentum transfers. Such experiments provide strong evidence for non-nucleonic phenomena. The purely nucleonic description fails badly for this isovector magnetic transition, which corresponds to the ${}^3S_1+{}^3D_1 \rightarrow {}^1S_0$ spin-flip transition. The OPE potential means the interchange of charged pions in the NN force. This is inevitably accompanied by 'transition radiation', since the light pion is suddenly accelerated and decelerated in the transmission from one nucleon to another. This is the origin of the so-called meson exchange contribution. As for the deuteron wavefunctions, a major problem has always been whether more complicated multi-pion exchanges are important or not. The whole meson exchange approach looks at first arbitrary, since there could in principle be uncontrollable contributions from complicated states of large numbers of mesons at short distances. Riska (1985) has pointed out that present empirical knowledge of the NN interaction is so good that it allows the explicit construction of the exchanged quantum numbers and their mass distribution in a nearly model-independent manner. This information can then be used to suppose that these objects are 'point like', and that they can then be coupled straightforwardly to the electromagnetic field. This so-called minimal coupling means that every time one has a momentum operator of a charged particle, it should be replaced by the usual relation $\boldsymbol{p} \rightarrow \boldsymbol{p} \pm e\boldsymbol{A}$, where $\boldsymbol{A}$ is the vector potential. When one applies this approach to the deuteron electro-disintegration at threshold, there is not only spectacular agreement with experiment without extra assumptions (Fig. 9), but it has also been found that it corresponds mainly to pion exchange (Mathiot 1984, 1985) The interaction densities for these transitions are insensitive to the uncontrollable short distance region (Fig. 10).

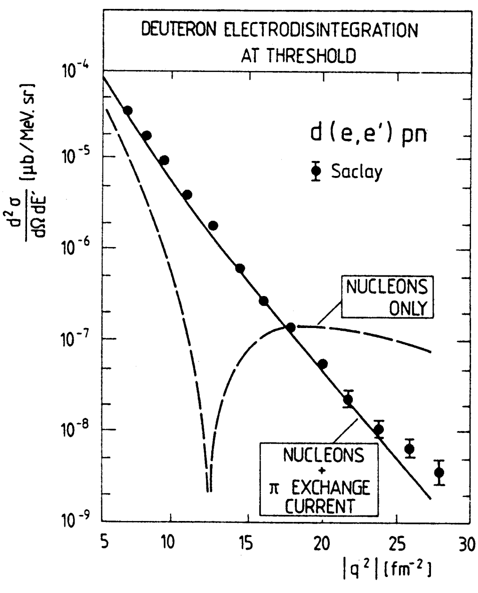

Fig. 9. Meson exchange effects in backward deuteron electro-disintegration near threshold, as a function of momentum transfer (Mathiot 1984).

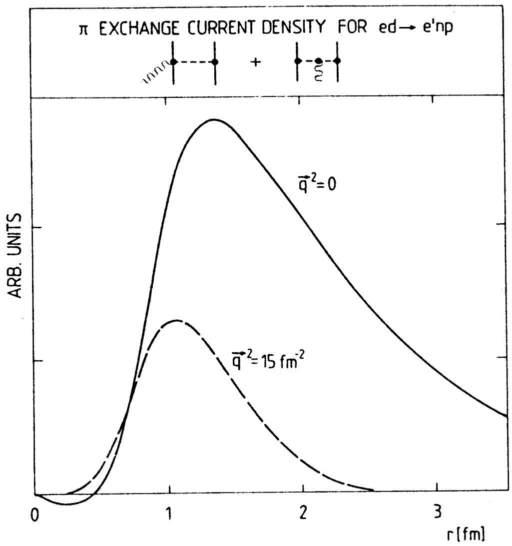

Fig. 10. The OPE transition density in the process ed → e′np at threshold (Mathiot 1985).

The upshot of all this is that there is an extremely good test of the validity of pion exchange in the case of the two-nucleon interaction, which is to say that we have a test of the major part of the nuclear force picture. The deuteron serves as an important constraint on any discussion of nuclear pion physics in more complex situations as well as on the more exotic quark physics. There are, however, additional contributions in nuclei, for example from ρ and Δ physics. The deuteron illustrates only the most important feature of meson physics, but it does so spectacularly. Do the dominant pion features show up clearly in other light nuclei? Indeed they do, and their signature is the tensor force which induces d-states in the $A = 3$ and $A = 4$ wavefunctions, which manifest themselves very clearly. The most striking case is that of the α particle, which is a manifestly deformed nucleus and not at all the splendidly symmetric object traditionally envisaged by most nuclear and particle physicists.

How does this deformation come about? The easy way (but not the quantitative one) is to visualise part of the α particle wavefunction as being due to two deuteron-like objects with their spins pointing the same way for a total spin of $S = 2$. An orbital $L = 2$ component must then compensate this spin so as to give a total $J = 0$ as schematically indicated in Fig. 11. It is this orbital component that is generated mainly by the OPEP. This deformation means that there are tensor correlations between the nucleons. In this way the α particle deformation becomes the prototype for those nuclear tensor correlations in heavy nuclei that are so important for an understanding of the magnetic moments and the axial current. We should look at the α particle as our smallest sample of nuclear matter. Experimentally this d-state has been discovered in various experiments with polarised deuterons. The most striking example is the d+d → ^{4}He+γ process. Since the deuterons are identical particles at low energy, the normally predominant E1 radiation is suppressed and E2 radiation dominates the threshold region. Owing to the d-state in the α particle, radiation can arise from an $L = 0$ initial state, beating the normal centrifugal suppression. This process occurs by the E2 $^5S_2 \rightarrow {^5D_0}$ transition, as seen for the astrophysical factor shown in Fig. 12. From such data one concludes that the α particle contains about 5–7% d-state (Assenbaum and Langanke 1987).

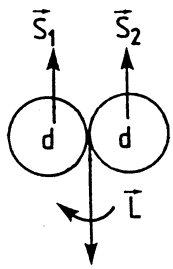

Fig. 11. Schematic picture of the ^{4}He d-state.

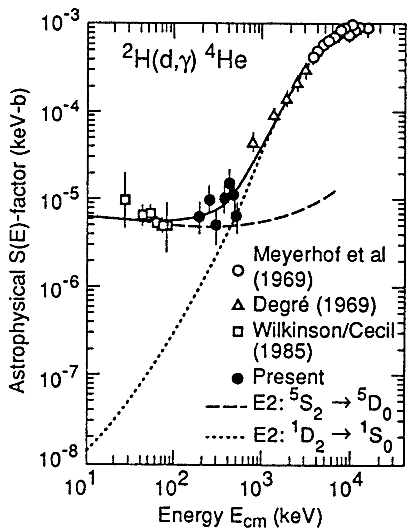

Fig. 12. The astrophysical factor for the reaction $^2H(d,\gamma)^4He$ (Barnes *et al.* 1987).

4. The Pionic Mode in the Nuclear Medium

I have presented a strong argument that the pion is essential to the dynamics and understanding of the light nuclei. Let us now turn to the properties of physical pions in the nuclear environment. The picture we have of a physical pion propagating through the nuclear medium is dominated by s and p wave scattering of the pion from the individual nucleons in the nucleus. In analogy with the p wave dipole scattering of light in a dielectric medium, the effects of pair correlations between the scatterers are crucial to a proper description: just as in a dielectric, the fact that scatterers cannot overlap leads to the notion that the field felt by an individual nucleon is an effective pion field and not the average field in the nuclear medium (Ericson and Ericson 1966). The code words for this are Lorentz–Lorenz effect, Landau–Migdal parameters, etc. One of the most influential contributions of pion physics has been to transmit this insight into nuclear many-body physics. These modifications are crucial, since it is only by such mechanisms that nuclear matter remains stable instead of reorganising its spin–isospin structure into a 'pion condensate', a reorganisation of the nuclear spin–isospin tensor correlations, which would take place if the nuclear tensor force were very strong. The consequences of such pair correlations appear for virtual pions in the nuclear many-body problem, for scattering of pions in the Δ isobar region, as well as in low energy pion–nuclear scattering. For the direct exploration and isolation of such effects, the energy region below 60 MeV kinetic energy is particularly attractive, since there the πN interaction is relatively weak.

Outside this threshold region the main driving force for pion propagation is the elementary p wave πN scattering amplitude averaged over spin and isospin, which can be schematically written

$$t_{\pi\mathrm{N}} = c_0 \boldsymbol{q} \cdot \boldsymbol{q}' .$$

It has exactly the same form as the long-wavelength limit of the scattering amplitude of light from an atom. As a consequence the potential or self-energy for pion propagation in the nuclear medium takes the form

$$\Pi = \nabla . \chi \nabla ,$$

where $\chi(r)$ is the pionic susceptibility of nuclear matter. It is, to leading order, given by the scattering strength c and the density $\rho(r)$ of scatterers:

$$\chi_0(r) = 4\pi c \rho(r) .$$

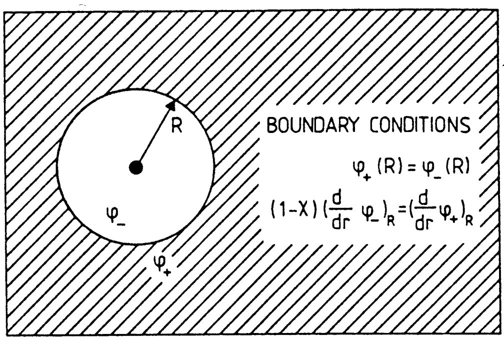

Fig. 13. Typical boundary conditions for the pion dipole field inside a cavity of polarisable nuclear matter.

This is an attractive interaction that is so strong that for normal nuclear matter it would overcome the repulsion from the kinetic energy and cause the collapse of the nucleus into a pion condensate. Matter exists because of the repulsion between nucleons that prevents any two of them to coexist at the same point. This causes a renormalisation of the interaction strength that is weakened to acceptable values. Fig. 13 displays the physical mechanism: it is exactly the same mechanism as in a dielectric. The p wave scattering leads to polarisation of the medium with induced (axial) dipoles appearing on the walls of a cavity and a reduced field on a test body in the centre measuring the effective field. This pionic Lorentz–Lorenz effect leads to the renormalised pionic susceptibility

$$\chi_0 = \chi_0/(1 + g'\chi_0) ,$$

where the classical value for g' is 1/3.

In a uniform medium this potential means that the energy–momentum relation for the pion has been modified from its free value

$$\omega^2 = k^2 + \mu^2 ,$$

where ω is the energy and k the momentum. In the medium this relation changes to

$$\omega^2 = (1 - \chi)k^2 + \mu^2 ,$$

from which we see that $\chi = 1$ is a critical point at which bound states with $\chi < \mu$ develop.

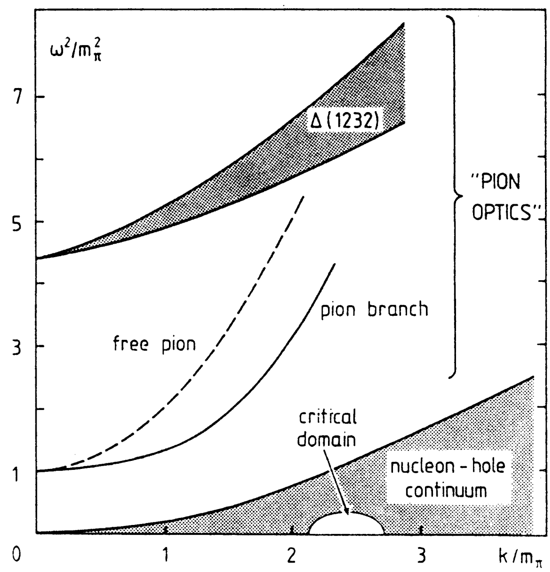

Fig. 14. The main branches of the spectrum of pion-like excitations in symmetric nuclear matter (Ericson and Weise 1988).

Fig. 14 gives the schematic energy–momentum relation for this pion branch in comparison with the free pion as well as the free Δ isobar. At small ω we see the nuclear spin–isospin excitations with pion quantum numbers and the critical region in which a pion condensate might occur. The most important feature of the pion branch is that it always lies below the free pion value. It is thus impossible to explore the nuclear pion branch using ordinary pion scattering as a probe. The question is then 'How can we display the properties of this collective nuclear mode?' The method that first comes to mind is to use the photon as a probe, since it penetrates deeply into the nucleus and it excites the Δ isobar strongly. This does not work. The photon is transverse, while the pionic mode is longitudinal, so there is no coupling. It is instructive to see how this difference comes about and what the consequences are. The Δ is a magnetic M1 excitation of the nucleon. If we then compare the coherent scattering amplitude from nucleons in the nucleus for π^0 propagation with the average neutral photoproduction amplitude, these have the following structures in terms of the incident momentum $\boldsymbol{k}$, the outgoing momentum $\boldsymbol{q}$ and the photon polarisation vector $\boldsymbol{\epsilon}$:

$$\begin{array}{ll} \pi^0 \text{ scattering:} & \text{constant} \cdot \boldsymbol{k} \,.\, \boldsymbol{q} \\ \gamma \to \pi^0 \text{ production:} & \text{constant} \cdot (\boldsymbol{\epsilon} \times \boldsymbol{k}) \,.\, \boldsymbol{q} \,. \end{array}$$

In a uniform medium the wave goes forward, so that $\boldsymbol{k}$ and $\boldsymbol{q}$ are colinear. This means that a π^0 wave produced by the photon vanishes: there is no coherent rescattering in this case as a consequence of the transversality. For π^0 scattering there is no such suppression, since the coherent wave is longitudinal. Fig. 15 shows that experiments bear this out. The total cross section for photons in the Δ region coincides closely with the average free cross section, which means that the Δ is virtually unaffected by rescattering in the transverse channel. In the case of pion scattering there are very strong changes in the longitudinal channel. This follows from the attenuation of the coherent pion wave in the medium.

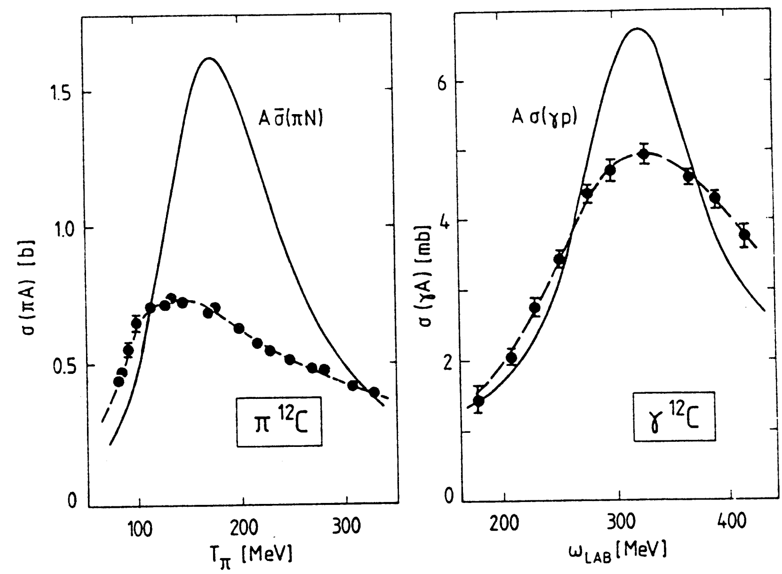

Fig. 15. The total π and γ cross section per nucleon versus energy for ^{12}C.

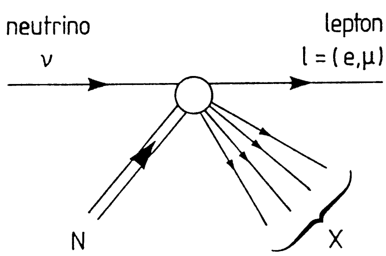

Fig. 16. Schematic picture of forward inelastic lepton scattering.

There is no effective probe available for the longitudinal (ω, k) response function (i.e. cross section). This is unlike the case for the transverse response, which can be explored by inelastic electron scattering. In principle an 'ideal' probe would be inelastic lepton scattering (ν, e or μ) near 0° as schematically indicated in Fig. 16. Here the interaction is longitudinal and one explores the energy–momentum relation $\omega = k$, which cuts right across the pion branch in Fig. 14. The response would be strongly enhanced at that point. The terrible price is of course the small cross sections and the poor energy resolution at present. We will now see how this problem can be solved.

5. Virtual Pion Beams and Pionic Neutrino Physics

Let us examine the neutrino reactions more closely. There is an instructive lesson that points to a way out of the dilemma using 'virtual pion beams', which can be achieved under particular conditions. For neutrinos this connection is the one between axial currents and pion physics, which can be given nearly rigorously within a consistent physical picture up to relatively large momentum transfers.

One of the important developments in nuclear physics in the last decade has been the discovery that forward np charge exchange reactions are a powerful tool to excite nuclear spin–isospin collective states, the Gamow–Teller mode. This line of research is now being pursued in producing collective modes at higher excitation in the Δ isobar region. We should not look at these excitations as being produced just by a nucleon–nucleon spin–isospin amplitude in multiple scattering: it is much more interesting and illuminating to consider that they are produced by a *beam of virtual pions.* We now have the prospect of exploiting such beams systematically. I will come back shortly to the physics we can learn from such beams in the np case, but let me first discuss what a virtual pion beam means and what the problems of using it are. To illustrate the nature of a virtual pion beam, consider the prototype case of 0° neutrino reactions in Fig. 16 above. In this case we can think of the incoming neutrino as decomposing virtually into a pion by the inverse pion β-decay $\nu \to \text{lepton} + \pi$, so that the incident neutrino is partially a cloud of pions, a beam of virtual pions. The consequences follow from a theorem. Its strong and precise predictions are nearly untested experimentally, unfortunately.

Adler's (1965) theorem: *Any 0° neutrino reaction with any final products on any target is exactly equivalent to the corresponding pion process with the energy–momentum of the pion given by the energy–momentum transfer ω, k, but the pion is virtual with $\omega^2 - k^2 = 0 \approx m_\pi{}^2$.*

The pion is virtual but the mismatch in momentum required to make it real is only about $\delta k \approx m_\pi{}^2/2\omega$, which becomes very small for large energy transfers. So the pion is virtual with novel physics but not so virtual that we completely lose contact with our usual physical picture of nuclear pion interactions. By the uncertainty principle it behaves as if it were a real pion over distances of $2\omega/m_\pi{}^2$, which easily attains nuclear dimensions. In the limit of large energy transfer, corresponding to very high incident energy, the virtual pion in 0° neutrino scattering will behave as a physical one! In this condition, the virtual pion mode joins with the physical pion branch in the dispersion curve.

This is an extraordinary result: the weakly interacting neutrino will mimic *all* the strong-interaction physics of the pion, albeit with the cross sections reduced by the weak coupling. The question is now whether we can achieve approximately this situation in practice at a lesser cost, but with the same physical insight. The answer is: yes. A proton surrounds itself by a pion cloud, since it is part of the time $p \rightarrow \pi^+ + n$, and similarly ^3He surrounds itself by a pion cloud through the virtual process $^3He \rightarrow {}^3H + \pi^+$. These virtual pions ride along with the incident proton or helium-3 and make a controllable beam. Indeed, it is easy to see from experiments that such forward processes are intimately linked to pion physics. Take for example the forward np → pn differential cross section in Fig. 17. The shape of the cross section does not depend on the incident energy when the momentum transfer is of the order of the pion mass or less. The typical scale of the dependence of the cross section on the momentum transfer follows from the pion mass.

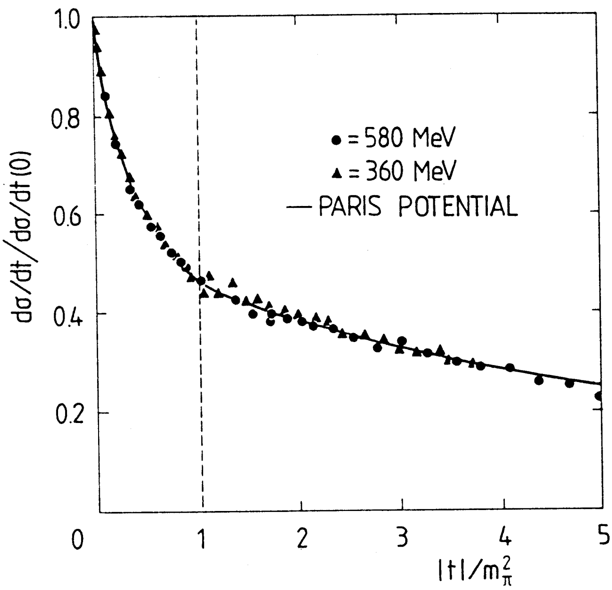

Fig. 17. The universal pion-dominated shape of forward np → pn charge exchange cross sections (Ericson and Weise 1988).

At strictly 0° the energy–momentum transfer will, to leading order, correspond to the energy–momentum of the effective pion in this virtual beam, as for neutrino reactions, but in the case of nuclei there will be complications due to distortion and multiple scattering and other perturbing effects which complicate the analysis: this is the price we have to pay for healthy cross sections. The question is whether we know how to control these effects. If not, such virtual beams will be without value in the pion perspective. And further: 'What can we now learn from this?' Consider the fundamental relation between excitation energy and momentum transfer for the spectrum of pion-like excitations in the nuclear medium illustrated in Fig. 14. The free pion has the relation $\omega^2 = m_\pi^2 + k^2$. The Δ isobar excitation for moderate momentum transfers in the nucleus has its strength located at higher excitation with its position, width and strength modified in the nuclear environment. At the bottom of the figure are the nucleon–hole continuum excitations, spin–isospin excitations, with the pion-like quantum numbers $J = 0^-$, 1^+, 2^-, etc. and $I = 1$. Below the free pion line is the 'pion branch', which reflects the modified energy–momentum relation of a pion propagating in a nuclear medium. The simplest way to visualise this mode is with the optical potential, but the mode as such is much more general. We expect this mode to fall below the free pion one, since we know that the effect of the repulsive kinetic energy is balanced by the attraction from the velocity-dependent π–nuclear potential. If we therefore think of this figure in

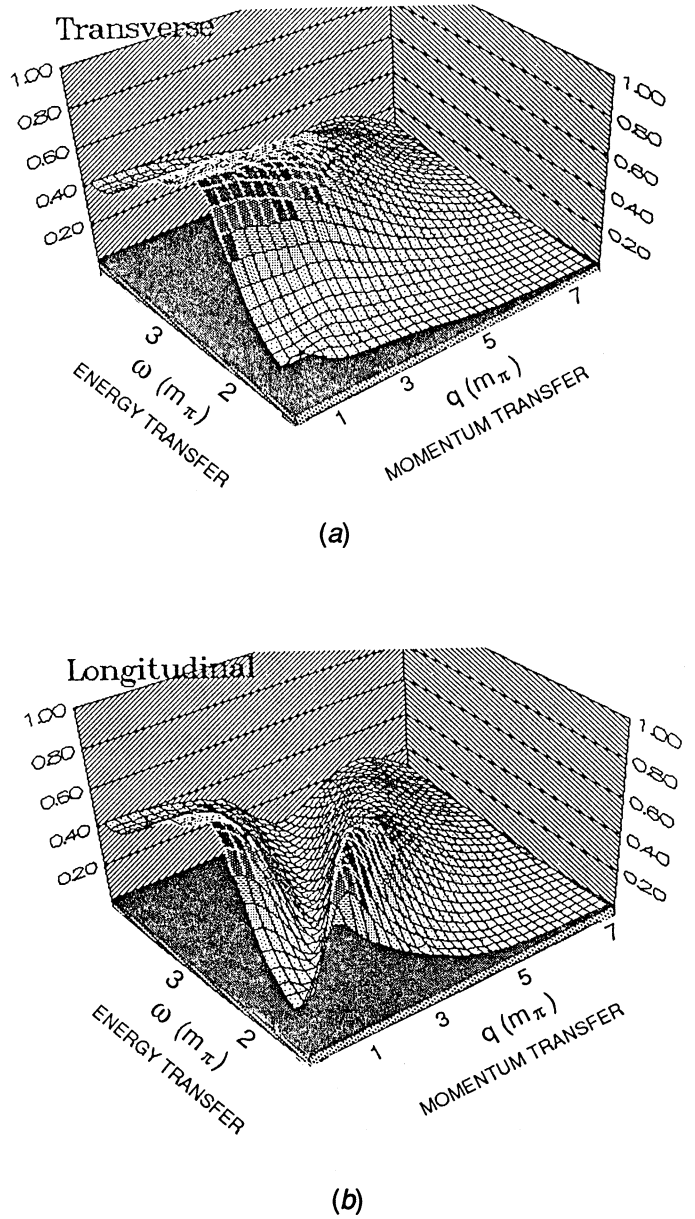

Fig. 18. The expected behaviour of the (a) transverse versus (b) longitudinal response surface in heavy nuclei (Delorme and Guichon 1989).

terms of a response function, it means nuclear matter will not react much to a pion-like perturbation with the free energy–momentum relation. The strength of the free pion branch on the one hand and of the nuclear Δ excitation on the other hand has been shifted into the pion branch with a new energy–momentum relation. The question is now whether we can observe this phenomenon, which is basic to nuclear pion physics. First, free pions give very limited information on the overall behaviour of the pion-like spectrum and its response function, although we do get information on the optical potential from elastic scattering: the exploration of these collective modes necessarily requires 'virtual' pions! As I discussed above, photons can in principle explore this, since they have $\omega = k$ as well, but the photon is transverse and does not have the pion quantum numbers. It does not couple to the pion branch. Instead, one observes beautifully that the Δ resonance survives nearly intact with nearly unchanged strength even in the heaviest nuclei. Thus, the exploration of the pion-like spectrum requires a longitudinal probe. The excitation landscape should then change dramatically as illustrated in Fig. 18.

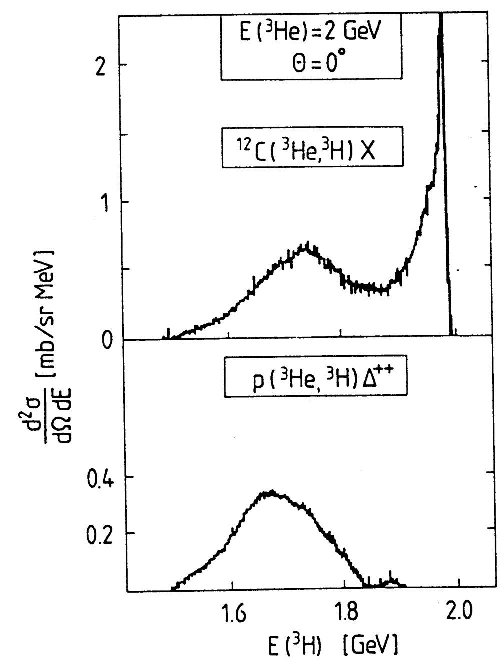

Fig. 19. The excitation spectrum for the energy loss in forward ^{3}He charge exchange on ^{12}C and on the proton.

If we now use the $^3\mathrm{He} \to {}^3\mathrm{H}$ forward charge exchange reaction, it is indeed possible to explore a part of this diagram outside the free pion kinematics. This was first pointed out by Chanfray and Ericson (1984). Fig. 19 gives a typical cross section for the process. It has a spectrum associated with spin-flip Gamow–Teller excitations at low energy, but in addition it shows a strong enhancement in the region below the Δ resonance energy (Ellegaard *et al.* 1985). It is the last feature that is important here.

Fig. 20 shows that there is a systematic 75 MeV downward shift in complex nuclei as compared with the leading-order 'free' response for the 'elementary process' on the proton. Even more striking, this shift is nearly the same both for ^{12}C and for ^{208}Pb! An analysis by Delorme and Guichon (1989) strongly suggests that we indeed deal with the explicit observation of the pion branch in nuclear matter and that various distortions and other perturbations can be handled quantitatively. They demonstrated that the triton spectra in Fig. 20 can be well understood both absolutely and as a function of angle, even though the cross section may vary by large amounts. The approach is based on a picture of π nuclear interactions, which contains a standard Δ isobar model with antisymmetry and isobar exchange, so as to give a consistent Landau–Migdal parameter g'. It was found that in practice the process divides about equally between longitudinal and transverse contributions. Of course, the pion shows up only in the longitudinal one, as in the neutrino processes. It seems reasonably safe to accept that we understand quantitatively this basic process to a considerable degree and that it has the required degree of sensitivity to pion physics to be a useful tool. The conclusion is that the observed shift has a major component from the collective pion mode in nuclear matter and that these experiments give concrete evidence for the existence of that mode.

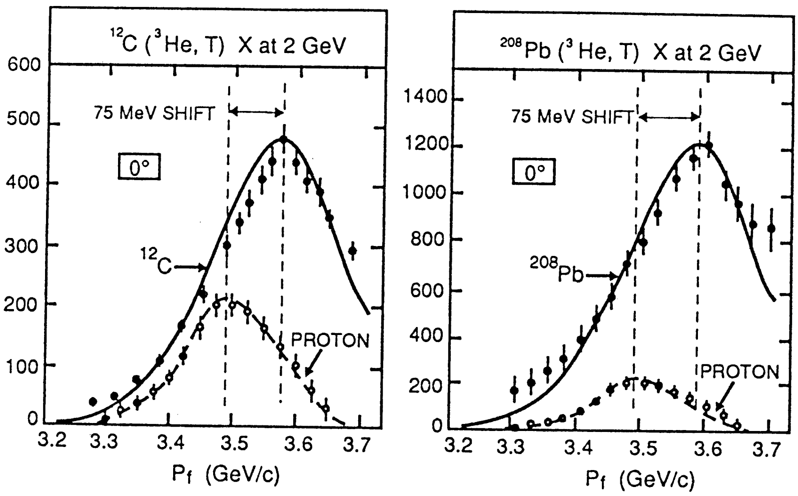

Fig. 20. Evidence for observation of the pionic mode in nuclei (Delorme and Guichon 1989).

This is the first clear evidence for this collective pionic mode, which differs profoundly from all previously observed collective nuclear states. One of the important tasks will be to find other probes and conditions that will display the pion branch optimally. The trickiest points at present have to do with the absorption and distortion of the projectile. This is why neutrinos would come in handy, but with formidable practical obstacles. In this case the kinematics are such that dramatic variations in the (ω, k) plane are expected even for small changes of $\omega = k$ in the region near 100 MeV pion kinetic energy (Delorme and Ericson 1985), as can be seen from Fig. 18.

6. Nuclear Chiral Symmetry and Spin–Isospin Physics

In quantum chromodynamics (QCD), chiral symmetry is broken by the small, but non-vanishing, masses of the u and d quarks. In the chiral limit these vanish and then the pion mass vanishes also. It is this aspect that is particularly important to nuclear physics. This soft limit, $m_\pi \to 0$, has a significance similar to that of the $\omega \to 0$ limit for photons, which leads to the universal Thomson amplitude for soft photons. In this case of long wavelength, the structure of the target becomes irrelevant and cannot be explored. The scattering depends then only on global quantities like the total charge and the total mass in the photon case. For nuclei, which are much larger than the pion Compton wavelength, such theorems remain valid only when scales are appropriately conserved. As was pointed out by Ericson *et al.* (1969, 1972), the pion Compton wavelength of 1·4 fm can be considered a large quantity, not on the level of the nuclear radius but on that of individual nucleons. These then appear structureless, so that the soft limit is true, but on the microscopic level only. The close link of the pion to axial phenomena appears in a wider perspective when seen considering chiral symmetry. In analogy with the photon that is associated with a conserved electromagnetic current and therefore with a conserved charge, one of the most important consequences of the nearly realised chiral symmetry is that the axial current A_μ is a nearly conserved quantity with a divergence proportional to the pion field ϕ [the partially conserved axial current (PCAC)]:

$$\partial_\mu A_\mu = -f_\pi m_\pi^2 \phi \, .$$

The consequences of chiral symmetry for nuclei follow nearly entirely from the validity of this relation. In this perspective we should think of the axial current not just as a spin–isospin operator but as an integral part of pion physics. In particular, this has the consequence that the individual nucleon spin–isospin operators for the axial current have induced terms due to pion rescattering as a natural complement. This is what we usually call pion exchange currents:

$$A_{\mathrm{a}} \propto \textstyle\sum \sigma_i \tau_l^{\mathrm{a}} + \text{exchange contributions:} \qquad \text{p wave pions,}$$
$$A_{\mathrm{a}}^0 \propto \textstyle\sum (\text{velocity terms and exchange contributions}): \quad \text{s wave pions.}$$

This result follows quite generally directly from PCAC and from the 'axial locality' hypothesis (Bernabéu and Ericson 1977). More precisely, the space-part of the axial current corresponds on the microscopic level locally to p wave pion interactions, and the time-part corresponds to s wave pion interactions, both on the level of individual nucleons. Therefore, the axial current operator acts exactly like an external s or p wave pion field on the nucleus. The pion source and the axial current not only have similarities but are equivalent in this limit.

One finds in this way a beautiful and remarkable analogy between the pion field and axial current in the nucleus and the potential and field in a dielectric. Ericson and her collaborators (Delorme *et al.* 1976; Ericson 1978) have emphasised and developed this picture.

This close parallel is no coincidence: in a dielectric we have permanent and induced electric dipoles, which react to the applied gradient of the potential; for the pion we have permanent axial dipoles (nucleon spins) and induced axial dipoles (p wave scattering) which react to the local gradient (i.e. p wave) of the pion field. This picture provides a powerful heuristic tool in our thinking. All kinds of dielectric phenomena have immediate analogues in the pion physics of nuclear axial currents. For instance, it becomes natural that the axial coupling constant g_A is renormalised in the nuclear medium by the Lorentz–Lorenz effect. Table 2 demonstrates that various quantities in the theory of dielectrics have an exact correspondence in nuclear pionics. It brings home the point made elsewhere at this Congress that different fields of physics have a unity in concepts and that we can learn much about one field by looking at another one.

Table 2. The analogy between dielectric media and pion physics

Dielectric quantities		Pionic quantities	
Potential	Φ	Pion field	φ_a
Displacement vector	$\boldsymbol{D}$	Axial current	$\boldsymbol{A}_a$
Divergence	$\nabla.\boldsymbol{D}=0$	PCAC	$\nabla.\boldsymbol{A}=-f_\pi m_\pi^2\varphi_a$ $(\rightarrow 0 \text{ for } m_\pi \rightarrow 0)$
Electric field	$\boldsymbol{E}=-\nabla\Phi$	Pionic axial current	$-\nabla\varphi_a$
Permanent dipole	$4\pi\boldsymbol{p}$	Nucleon axial dipole	$-(g/m)\sigma\tau_a$
Induced polarisation	$4\pi\boldsymbol{P}$	Axial exchange current	$\boldsymbol{A}_a^{ex}$
Dielectric constant	ϵ	Axial susceptibility	$(1-\chi)$

7. Conclusion

Explicit and quantitative test cases now display the dominant role of the pion on the microscopic level in nuclei. Here the deuteron provides the most spectacular example. This transmits itself into important nuclear properties, not only into the binding energy, but in particular in the form of the nuclear tensor correlations. The d-state of the deuteron and the recently observed deformation of the α particle are striking and specific illustrations. In the discussion of pion physics in nuclei we are also constantly confronted with the issue of nucleon pair correlations and their effects. These are responsible for the effective pion fields in nuclei. The explicit observation of such phenomena in double charge exchange reactions, where they can be clearly displayed, strengthens our confidence in this description of the pion in the nuclear medium. A major advance in our understanding is the deep relation between the axial spin–isospin phenomena and nuclear pion physics. This leads to a detailed interpretation of the axial current which closely parallels the theory of the Maxwell displacement vector and the susceptibility of an inhomogeneous dielectric. The main consequences are contained in the PCAC and axial locality relations. We have now a consistent physical framework within which we can interpret the pion-like response function and its physics. The recent explicit observation of the pion collective mode is a novel feature of the nuclear many-body system. The ideal tool for exploring its properties and other aspects of axial pion physics in nuclei would be virtual pion beams provided by neutrinos. Although untraditional to nuclear physics, this is an area which should be closely watched. The relevant experiments could be realised with the high-intensity kaon factory KAON that very likely will be built in Vancouver as an extension of the TRIUMF meson factory.

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