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NUCLEAR PIONIC MODES : POLARIZATION MEASUREMENTS, THE CLUE FOR A DEFINITIVE IDENTIFICATION (1)

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Résumé - On présente les progrès théoriques récents dans l'interprétation des réactions d'échange de charge sur les noyaux dans la région de la résonance Δ. Le principal résultat est l'attribution au mode nucléaire pionique d'une bonne partie du décalage d'énergie vers le bas observé par la réaction (\(^6\)He,t). Une condition non encore vérifiée est que le canal d'excitation pionique soit suffisamment favorisé dans la réaction. On discute comment une mesure d'observables de polarisation peut permettre une identification définitive.

Abstract - Recent theoretical progress in the interpretation of charge exchange reactions on nuclei in the Δ resonance region is reviewed. The main outcome is that an important fraction of the downward energy shift observed in the (\(^6\)He,t) reaction is likely to be attributable to the nuclear pionic mode. A condition not yet proven however is that the driving interaction goes appreciably through the pion-like channel. It is discussed how measurement of spin observables can help to make a definitive decision.

1. INTRODUCTION

Measurements of the spin responses of nuclei at low frequencies (< r-mass) have an already long and deceptive story. Along years much excitement has been invested in the quest of pion condensates and their precursor phenomena. The most crucial experiments have been those involving polarization since they permit in principle the separation of the spin-longitudinal or pion-like response (to be denoted hereafter by SL) where the interesting effects are expected from the transversal or photon-like one (ST). They have practically given the death-blow to the hope of observing the last relics of pion induced spin order in the nuclear medium/1/. Though the situation is not completely clear there is a general belief that the short range repulsion is strong enough to kill all manifestations of the r-induced attraction in the spin-isospin channel. Present progress is made in the direction of separation of the spin isoscalar responses where some exciting predictions of theory will hopefully have a more fortunate fate.

The situation changes drastically beyond the pion threshold where very positive signals of collective phenomena generated by pion propagation have been gathered by systematic studies of charge exchange reactions mainly conducted at the Saturne accelerator. The main feature revealed by those measurements of the nuclear response function in the Δ resonance region is the existence of a universal downward shift of 60 – 80 MeV of the resonance position compared to that measured on a proton target/2-4/. This observation was soon taken/5/ as a signature of the pionic mode or pionic branch, a high lying collective mode of the nucleus/6/. After a number of converse statements/T-10/, elaborate calculations have confirmed that the positive conclusions of ref. /5/ were qualitatively right though only a fraction (= 30 MeV) of the shift was found to originate

* The content of this talk owes much to work made in collaboration with P. Guichon and P. Desgrolard.
in the collective effect/11,12/*/ its larger part arises from the deformation of the \( \Delta \) peak in the nucleus (by Fermi motion, Pauli blocking of the \( \pi \)-nucleon decay and coupling to reaction channels) as viewed by a momentum transfer dependent probing interaction including the form factors at the projectile-ejectile vertex.

The nearly parameter-free calculation to be presented here is now in excellent agreement with experiment in the most documented case of the \( (^3\text{He},t) \) reaction at 2 GeV. Nevertheless a weak point remains due to uncertainties in the elementary driving interaction \( NN \rightarrow N\Delta \) which has been constructed from a \( \pi \) and \( \rho \) mesons exchange potential. It is indeed necessary that it contains a sizable \( SL \) component since a \( ST \) excitation would be essentially blind to the nuclear pionic mode as it is for other pion propagation effects/13/.

Measurements of polarization observables are thus necessary to ascertain the longitudinal/transversal ratio.

This paper is organized as follows. The salient features revealed by measurements of the nuclear response around the \( \Delta \) resonance are first briefly reviewed with comparison of the results obtained by various probes. The theoretical ingredients for calculation of the response with full account of pionic correlations are summarized. For the sake of illustration a schematic two-levels model is presented which is shown to reproduce qualitatively the main features of the realistic calculations (a more extended account of our work can be found in ref. /11/ and a forthcoming publication).

The computational steps from the nuclear response to the charge exchange cross-section are described and the results compared to experiment. Finally the open problems are discussed with special emphasis on the crucial informations which are or will be brought in by measurements of polarization observables.

2 - DELTA'S IN NUCLEI: A SURVEY OF THE EXPERIMENTAL SITUATION

Numerous reactions permit excitation of the \( \Delta \) in nuclei and many have actually been used with the aim of seeking deviations of the resonance properties in the nuclear medium compared to those measured in free space. It is convenient for our purpose to put the emphasis among the gathered informations on the modification of the resonance position (the shift) and the scaling law vs. mass number \( A \) of the response on top of resonance. This is certainly an over-schematization of the available observations. For instance we could have equally retain the evidence for the modification of the width since as will appear below it plays a very important role in charge exchange reactions through generation of a shift which has nothing to do with pionic correlations.

2.1 - \( \gamma \) and \( \pi \) probes

The two selected features are summarized in Fig. 1 for probes which couple to the \( N\Delta \) excitation through pion or photon exchange, including the case where these particles are real beam constituents. Each time we have distinguished the \( SL \) or \( ST \) nature of the coupling according to the respective projections \( S \cdot q \) or \( S \times q \) of the usual \( 1/2 \rightarrow 3/2 \) spin transition operator. It appears that charge exchange of proton and light ions share some properties akin to both \( \gamma \) and \( \pi \) excitation. Like the former they display a universal behaviour for all targets (with of course the exception of the scaling law). Nevertheless a strong downward shift is observed like in the \( \pi \)-nucleus total cross-section. How these observations can be reconciled is explained in the next two sections.

3 - CALCULATION OF THE RESPONSE

The nuclear response at given momentum-energy transfer \( (q,\omega) \) is related to the full \( \Delta \)-hole propagator by:

\[
R(q,\omega) = -\frac{1}{\pi} \Im \Pi(q,\omega)
\]

Fig. 1 - \( \Delta \) properties in nuclei as seen by various probes.
Medium effects enter at two levels in the evaluation of this quantity (Fig. 2). One has first (1st term in r.h.s. of Fig. 2) to consider the renormalization of the particle (i.e. $\Delta$) lines (we have assumed that insertions in the nucleon hole lines play a minor role in the present problem). We have been content to include modifications of the $\Delta$ width i.e. with account for Pauli blocking and distortion of the $\pi + N$ decay channel and opening of 2 and 3 nucleons emission channels according to available parametrizations/14/. We needed no real part for the $\Delta$ self-energy when checking our calculations on $\pi$ and $\gamma$ measured cross-section. Anyway it is known that theoretical estimates of this quantity are not very reliable and that it can hardly be separated in phenomenological determinations from the effects due to short range $\Delta$-hole correlations. The zero-th order $\Delta$-hole propagator is easily obtained in nuclear matter with this procedure. It can be used as the basis of a semi-classical approximation for the case of finite nuclei since the large width of the $\Delta$ implies a smearing of the response over many shells and the washing out of quantum effects. We have thus evaluated the propagator as the first term of a Wigner-Kirkwood expansion for the two instances of $^{12}$C and $^{208}$Pb nuclei:

$$\Pi_0(q, q'; \omega) = \int d\rho \, e^{i(q+q')\cdot\rho} \, \Pi_0^{NM}(\frac{q + q'}{2}; \omega; p(p))$$

where $\Pi_0^{NM}$ stands for the propagator in nuclear matter evaluated for a position dependent Fermi momentum $p(p)$ determined from the local density $\rho(p)$. At this stage Fermi motion is thus taken into account.

The interaction between $\Delta$-hole excitations generates (spin-isospin) correlations which are most commonly treated within the RPA approximation. Its prominent component is pion exchange (OPE) which can be on its mass shell between two successive excitations since we are at excitation energies beyond the pion threshold. Dramatic effects can thus be expected from the crossing of the pole. Iteration of the r.h.s. of Fig. 2 generates chains of excitations where $\Delta$-hole and pion alternately propagate, successively converting in each other. There are several differences with more familiar collective phenomena at low excitation like the strong frequency dependence of the interaction, the large width and the overlapping of excitations of many multipolarities. Nevertheless a major property of collectivity is still present. The (free) pion dispersion line $\omega_i(q) = \sqrt{q^2 + m_i^2}$ separates the $(q; \omega)$ plane in two zones according to the sign of the OPE interaction: the response or strength is repelled upwards or attracted downwards according to the location of the considered point beyond or below the line.

The same phenomena appear in the optical potential approach which is often used for the scattering of (real) pions. Summation of the multiple scattering series plays the role of the RPA. As an introduction to the presentation of the results of the actual calculation the schematic two-levels model due to P. Guichon/15/ is especially illustrative. One starts from a unique zero-width $\Delta$-hole excitation at energy $\omega_\Delta(q) = \sqrt{M_\Delta^2 + q^2} - M_N$ (with $M_\Delta$ and $M_N$ the $\Delta$ and nucleon masses) and an uncoupled pion. The momentum $q$ plays the role of a parameter. The $\Delta$-hole propagator can be written as $\Pi_0(q; \omega) = 2\rho \frac{\omega_\Delta(q)}{\omega_\Delta^2(q) + i\epsilon}$ with $\rho$ the density of the medium. At
the point \((q \approx 2.1m_\pi; \omega \approx 2.4m_\pi)\) the two dispersion lines cross each other. Coupling the two levels through the \(p\)-wave \(\pi N \Delta\) vertex one gets the new \(\Delta\)-hole propagator:

\[
\Pi(q;\omega) = \Pi_0(q;\omega)/[1 - \Pi_0(q;\omega)^2/(\omega^2 - \omega^2_\Delta(q) + i\epsilon)]
\]  

(3)

A similar expression holds for the new pion propagator. Both have now two poles at \(\omega_\Delta^2(q,q) = \omega^2_\Delta(q) + \omega^2_\pi(q) = \sqrt{2(\beta^2_\Delta(q) + \beta^2_\pi(q))} ± \sqrt{2(\beta^2_\Delta(q) + \beta^2_\pi(q))^2 + 8\omega^2_\Delta(q)q^2}/2\). Hence there appear two modes which no longer cross, the delta branch at energy \(\omega_\Delta\) and the pion branch at the lower energy \(\omega_\pi\). These names coming from the fact that they rejoin the \(\Delta\)-hole and pion lines at \(q = 0\) where the coupling vanishes. At the place of previous crossing they have become a superposition of \(\Delta\)-hole and pion with about equal amplitudes illustrating our previous discussion in terms of successive conversion and propagation of the coupled excitations. We have plotted the previous and new dispersion lines on Fig. 3 superimposed on the result of the realistic calculation so as to show that the qualitative features of the model are preserved. This quasi-particle aspect (hence the term quasi-pion sometimes used for the lower energy mode) has been especially emphasised in ref. [16].

The propagator (3) can be decomposed according to the two modes:

\[
\Pi(q;\omega) = 2\epsilon_\omega \omega_\Delta(q)[Z_+(q,\omega)/[\omega^2 - \omega^2_\Delta(q) + i\epsilon] + Z_- (q,\omega)/[\omega^2 - \omega^2_\Delta(q) + i\epsilon]]
\]  

(4)

where \(Z_\pm (q,\omega) = [1 ± \sqrt{1 + 8\epsilon_\omega(q)/(\omega_\Delta(q) - \omega^2(q))^2}/2\) such that \(Z_+ + Z_- = 1\). The \(\Delta\)-hole strength has been splitted in two parts which read:

\[
R_\pm(q;\omega) = \epsilon_\omega \omega_\Delta(q)[Z_\pm (q,\omega)/\omega_\Delta(q)]\delta(\omega - \omega_\pm)
\]  

(5)

Only the lower energy one, i.e. the pion branch, lies mostly in the space-like region and is easily accessible to experiment.

Taking into account the OPE only is an oversimplification of the interaction. A short range piece should be added characterized by a parameter \(g_A^\prime\) which is the analog of the Landau-Migdal parameter \(g^\prime\) used at low frequency. Its value is not well determined since as mentioned above it is hardly distinguishable from self-energy insertions in the \(\Delta\) lines. It commonly ranges between 0.33 and 0.50, the latter number being our working choice. It is far from playing the same role as its low energy analog \(g^\prime\) since the crossing of the pole makes the OPE to overwhelm the \(\Delta\)-hole interaction. It projects equally on the \(SL\) and \(ST\) channels. We have added in the latter a \(p\) meson exchange interaction which compensates somewhat the repulsion generated by \(g_A^\prime\). Anyway there is no crossing of levels in this channel so that nothing spectacular can be expected.

**Fig. 3** - Map of the response isolines in the \((q;\omega)\) plane. From left to right \(R_0, R_T, R_L\) (see text for definition). The dot-dashed lines are the free \(\Delta\) and \(\pi\) lines, the full lines on the right are the \(\Delta\) and \(\pi\) branches of the schematic model. For completeness we have also represented the photon (\(\gamma\)) and free nucleon (\(N\)) lines.

The realistic account of the RPA correlations is obtained by solving coupled (i.e. SL vs. ST) integral equations for each multipolarity \(J\) of the \(\Delta\)-hole propagator:

\[
\Pi^J(q,q';\omega) = \Pi^J_0(q,q';\omega) + (2\pi)^{-3} \int d^3q d^3q'' \Pi^J_0(q,q'';\omega) V_\Delta(q'',\omega) \Pi^J(q'',q';\omega)
\]  

(6)
These equations are of course singular in the case of unnatural parity where \( V_{A}(q^2;\omega) \) contains the OPE. Selecting our results for the case of \(^{208}\text{Pb}\) we display on Fig. 3 three results: \( R_{0}(q;\omega) \), i.e. in absence of RPA correlations, \( R_{2}(q;\omega) \) and \( R_{T}(q;\omega) \) according to the type of excitation operator at the two vertices of the full A-hole propagator. As expected there are no drastic modifications for the ST channel. Indeed the measured \( \gamma \)-nucleus total cross-section scales like the mass number \( A \). On the opposite there appears a complete remoulding of the landscape in the SL channel. The two region where A-hole strength accumulates are separated by a deep valley which is roughly followed by the real pion line. Hence we get the decrease in \( A^{-1/3} \) of the \( \pi \)-nucleus total cross-section per nucleon which is up to a factor the SL response along this line and a downward shift increasing rapidly with \( A \) of the saddle point which constitutes its maximum. As a corollary, exploration of the interesting regions in the \((q;\omega)\) plane means the use of virtual pions. Neutrinos would be quite an ideal source but are not very manageable. Baryon charge exchange reactions are an ersatz more easily available but as will be seen below have the drawbacks of a not too well known driving interaction and poor penetration in the nucleus.

4 - FROM THE RESPONSES TO CHARGE EXCHANGE CROSS-SECTIONS

Direct connection between the nuclear response and baryon charge exchange implies that a single step mechanism is valid which means that the projectile-ejectile is acting like a "lepton" probe at least with a sufficient accuracy. In particular one should insure that the excitation of the \( \Delta \) in the projectile which has been invoked as a source of the observed shift/9/ does not occur with much probability. Experiments are dearly reinsuring on this point. Indeed there has been search for \( \Delta \) excitation by the \((^{14}\text{N},^{14}\text{O})\) reaction with no success (see ref. /2b/). Due to nuclear structure this transition is much hindered and prevents one step excitation with appreciable amplitude. A two steps process with \( \Delta \)-nucleus intermediate states does not suffer from structure forbiddenness and would be observed if the two steps mechanism were an important contributor.

As represented on Fig. 4 the calculation factorizes in three steps in the plane wave approximation: 1) the nuclear response treated in §3, 2) the driving effective interaction \( V_{eff}(q;\omega) \) which plays the role of the \( \gamma \) and \( W \) bosons in leptonic interactions, 3) the form factors at the projectile-ejectile vertex where the structure of these nuclei enter.

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As concerns point 3) we specialize the discussion in this section to the case of the \((^{6}\text{He},t)\) reaction which is the most documented from both the experimental and theoretical points of view though
progress is currently done for other reactions. It is clear that two different form factors are needed according to the \( SL \) or \( ST \) nature of the \(^3\)He-\( t \) vertices. If the latter are restricted to one-body operators one expects a slower fall-off of the \( SL \) form factor vs. the \( ST \) one because S-D wave interference is constructive instead of destructive. Experimental information is unfortunately limited to the transversal case through magnetic electron scattering. In order to palliate our ignorance on the elusive longitudinal vertex we have first used a phenomenological strategy, with the \(^3\)He-\( t \) reaction on hydrogen target serving as calibration. We have thus extracted a gaussian parametrization in the squared \( t \) momentum \( q^2-W' \) for the \( ST \) form factor from the isovector combination of magnetic data obtained on \(^3\)He and tritium. The experimental spectra obtained on proton/\(^3\)He at 2 GeV have then been fitted to our theory with two free parameters, the slope \( \alpha_L \) of the \( SL \) gaussian and an overall normalisation factor \( N \). As was hoped the adjusted value 0.143 \( m_g^{-2} \) of \( \alpha_L \) turned out to be significantly smaller than the experimental number 0.230 \( m_g^{-2} \) for \( \alpha_T \). More surprising and probably with some luck the normalisation was found very close to 1 : \( N = 0.88 \).

![Fig. 5](image)

**Fig. 5** - The invariant cross-section \( d^2\sigma/dt ds \) of the reactions \( p^3\)He-\( t \)\( \Delta^+ \) (left) and \( p(d,3p)\Delta^0 \) (right) at \( \sqrt{s} = 1192 \) and 1200 MeV respectively for 2 GeV projectile energy. See text for the meaning of the curves labelled \( L \) and \( T \). Units are \( mb/GeV^4 \).

As an alternative to get rid of part of the preceding phenomenology one can undertake the full computation of the \( SL \) and \( ST \) form factors with Faddeev wave functions for the tri-nucleon system and one- and two-body operators at the two vertices (Fig. 4 right). We have achieved this goal/19/ with careful checking that we are able to reproduce the measured magnetic form factors and also the magnetic moments and Gamow-Teller matrix elements which serve as normalisation points for the \( T \) and \( L \) cases *. We have confirmed to a large extent the results of the phenomenological approach. This method has not yet been applied to the calculation of the triton spectra from complex nuclear targets. There is little doubt that it will produce results very similar to the previous one in view of the agreement found for the case of hydrogen target. To illustrate our claim we present on Fig. 5 (left) our no-parameter calculation of the invariant cross-section on top of the resonance as a function of \( t \) for the reaction on the proton. The comparison with experiment is quite eloquent (even the normalisation turns out to be right!). The curves labelled \( L \) and \( T \) represent the situation where both the \( SL \) and \( ST \) vertices have the same \( L \) or \( T \) form factors. It is clearly seen that the fall-off of the cross-section gives a hint for longitudinal dominance, a favourable situation for sensitivity to the pionic modes. However this is only an indication. What fixes the relative importance of the \( L \) channel is actually the ratio \( W_{L/T} = W_L/(W_L + 2W_T) \) where the weighing factors \( W_{L,T} \) are the products (form factor) \( \times \) (effective interaction) and the factor 2 accounts for the two transversal degrees of freedom. Our calculation shows a rapid increase of this ratio with increasing energy transfer. This is due to the convergence of two effects : 1) as already seen the \( SL \) form factor drops less rapidly than the \( ST \) one; 2) the \( L \) component of \( V_{eff} \)

* We have used the Reid soft core potential and two-body operators from \( \pi \) and \( \rho \) exchange
is dominated by the rise of the $\pi$-exchange whereas the $T$ component is rapidly decreasing due to cancellation of the short range repulsion by the $\rho$-exchange. As a result though the $L$ channel has the initial handicap of only one degree of freedom it overcomes the $T$ channel across the resonance with a $W_{L,T}$ value of 1/2 attained at about 275 MeV in the phenomenological approach and 295 MeV in the full theoretical one. On the right part of the figure are plotted the results of the similar calculation in the case of the $(d,2p)$ reaction. Though the first points are in good agreement with experiment/2c/ there is clearly a problem that will be discussed in more details in next section.

The last step to the computation of triton spectra is an evaluation of the distortion. At the relevant energies the absorption of the projectile and ejectile in the target is the prominent feature. We have made use of an eikonal approximation with single convolution of the $^3$He and tritium densities with that of the nucleus. The number of active nucleons which is only 1.6 in $^{12}$C and 3.4 in $^{208}$Pb gives an idea of the peripheral character of the reaction.

Fig. 6 - Comparison of theory vs. experiment for $^{12}$C $(^3$He,t) $X$ at $0^\circ$ and 2 GeV for $^{12}$C (left) and $^{208}$Pb (right). Units are $\mu b/(sr \times MeV/c)$ for the cross-section and GeV/c for the triton momentum $P_t$. See text for explanation of the curves.

Putting all ingredients together we have computed spectra at $0^\circ$ and $3^\circ$. Only the first angle is presented on Fig. 6 though the second case is also quite satisfying. We have applied normalisation factors of 0.75 and 0.80 respectively to $^{12}$C and $^{208}$Pb targets. The spectra for the proton target are also plotted in order to display the energy shift which amounts to 74 MeV in $^{12}$C and 83 MeV in $^{208}$Pb. Agreement with experiment/2a/ is extremely good except for the dip region (extreme right of the curve) which is generally poorly understood even for the cleanest probes. However comparison between the dot-dashed (only Fermi motion and $\Delta$ width modifications) and continuous (full calculation including RPA correlations) curves reveals that a major part of the shift would be present without spin correlations. The displacement of the $\Delta$-hole strength to the pionic branch adds only 25 to 30 MeV from $^{12}$C to $^{208}$Pb.

Understanding the origin of the non collective part of the shift that we will call a deformed shape induced shift (to be abbreviated as d.s.i.) is a rather subtle matter. We have seen in the beginning of this section that the actual reactions are not a direct measure of the response. The latter is weighed by the factors $W_{L,T}$ defined before. All are model dependent quantities which cannot be reliably disentangled contrary to the case of $(e,e')$ scattering. As will become clear below it is important to know how these factors vary across the resonance. In our model they have contrasted behaviours according to the spin channel. In the $SL$ one they first increase up to an energy transfer of 276 MeV because the rise of $OPE$ is not yet compensated by the slow fall-off of the $SL$ form factor and then decrease. In contradistinction the variation of both the $ST$ form factor and the driving interaction (the rise of $\rho$ exchange cancelling progressively the repulsion) conspire to a much faster decrease. It is easy to understand that the maxima of the spectra arise at higher energy than the maxima of the responses if the $W_{L,T}$ factors increase with increasing energy transfer and
lower if they decrease. Beyond 275 MeV excitation one gets thus downward displacements in both channels though much more pronounced in the ST one. Hence the cross-section $S_p$ on the proton has its maximum displaced downwards by 35 MeV (15 and 60 MeV for SL and ST respectively) with respect to that of the Breit-Wigner response. This is the point which serves as the reference for defining the shift (maximum of the dashed curve of Fig. 6). Of course the spectra $S_p$, $S_L$ and $S_T$ on nuclear targets undergo such displacements of their maxima with respect to those of the responses $R_p$, $R_L$ and $R_T$, the precise value depending on their detailed shape. The shift with respect to the proton data is then a differential effect between the respective displacements. A d.s.i. downward shift occurs whenever the maxima of the nuclear spectra undergo displacements larger than 15 MeV (L) and 60 MeV (T). As a first instance we consider the case of the RPA responses $R_p$ and $R_L$ of $^{12}$C which will constitute our test nucleus for this discussion. The correlations yield a downward collective shift of 65 and 23 MeV respectively with respect to the bare response $R_0$ (the latter number signing a slight overcoming of the p induced attraction over the short range repulsion). After due consideration of the $W_{L,T}$ factors one gets an extra downward d.s.i. shift of 44 MeV for the total spectrum which takes its source in values of 80 MeV for $S_T$ and 24 MeV only for $S_L$. The maximum of the latter is not easy movable not only because the decrease of $W_L$ is slow but also because the response $R_L$ is quite steep (cf. Fig. 3). Subtracting the corresponding numbers for the proton the total shift amounts to 75 MeV (about 78 MeV for $S_L$ and 50 MeV for $S_T$).

Another extreme attitude would be to ignore the existence of spin-isospin correlations. One would start then from the response $R_0$ which contains the non collective medium effects described in §3. Altogether they produce considerable broadening and flattening of the resonance along the ($^3$He, t) kinematical line as compared to the free proton (or neutron) situation. This reshaping is not specific to any spin channel and is clearly seen in ($e,e'$) data/20/. As a result, the flatter bare spectrum $S_0$ is moved downwards by as much as 80 MeV (60 and 96 MeV for $S_{L}$ and $S_{T}$) with a resulting d.s.i. shift of $\approx 50$ MeV after readjustment on the proton (60 and 40 MeV for the L and T bare spectra). The difference of 25 MeV with the case where RPA correlations are fully accounted for corresponds to the quantity which is read on Fig. 6 between the maxima of the full and dot-dashed curves.

In the preceding discussion we have actually been referring to the ideal case of a plane wave probe. Taking distortion into account we have to distinguish between central and peripheral multiplicities. The first ones have a large collective shift mainly in the SL channel with a maximum which can hardly be moved because of its location below 275 MeV and their steep rise. As for the second they are less sensitive to correlations and have thus a smoother shape with an easily movable maximum. Therefore they display an important d.s.i. shift. As a result of this subtle interplay the contributing multipoles have all more or less the same shift. The peripheral waves which remain in presence of absorption of the probe (the window is $J = 4-7$ in $^{12}$C and 10-15 in $^{208}$Pb) have as a mean 1/3 collective vs. 2/3 d.s.i. shifts. An indication of the role played by the form factors and the driving interaction to generate a large d.s.i. shift can be found in the decrease of the observed shift at higher bombarding energy/3/. If we assume that our model remains valid there, indeed the range of variation of the squared momentum transfer $t$ across the resonance then diminishes so that the non collective shift tends to disappear in the high energy limit.

5 - THE OPEN PROBLEMS AND THE ROLE OF POLARIZATION

The existence of a downward shift of the $\Delta$ resonance is not peculiar to the ($^3$He, t) reaction. This phenomenon is common to many projectile-ejectile couples. The prototype (p,n) reaction at 800 MeV displays a shift/4/ somewhat smaller than with the trinucleon probe (51 and 57 MeV for $^{12}$C and $^{208}$Pb). We interpret this difference by the absence of nuclear form factors. Indeed in our calculation the d.s.i. shift is much smaller than before for the RPA response. Our total shift amounts to 32 MeV only. A more serious drawback is that the resonance position turns out sensibly too low both for the proton and nuclear targets. This can be due only to some defects of our driving interaction $V_{eff}$. A careful analysis shows that it comes certainly from the antisymmetrisation procedure which we used to generate its short range part. There is an inherent dependence on the energy of the incident probing particle which causes a concomitant decrease of the repulsion with the energy. Therefore at 800 MeV we get a rapidly varying transversal component of $V_{eff}$ due to cancellation of $\rho$ exchange across the resonance. This is a feature which is not supported by polarisation data on the (p,n)$\Delta^{++}$ reaction which shows the presence of a strong ST component/21/. This does not necessarily invalidate our calculation for the ($^3$He, t) reaction where the energy per nucleon was sensibly lower so that we could generate just the right
amount of repulsion. It is clear however that this point has to be clarified by separation of the SL and ST contributions in view of the role of the cancellation occurring in the ST channel for the generation of large d.s.t. shifts. This is possible through the use of polarization which is rapidly developing.

There are already available data from the $(d,2p)$ reaction with a polarized beam/22/. Though a downward shift of 65 MeV is again displayed in $^{12}$C they indicate a dominance of the ST component which is even increasing from proton to $^{12}$C targets with respective $W_{L/T}$ ratios of 0.89 and 0.33 respectively. This may sign again a change of $V_{eff}$ at increasing energy (here 1 GeV per nucleon). Our effective interaction is completely inconsistent with such dominance because it has an even weaker ST component than for 800 MeV protons. How to reconcile both observations is a puzzling problem on which more theoretical and experimental work is needed.

![Angular Distribution](image)

**Fig. 7** - The angular distribution of the decay pions with respect to the direction of the momentum transfer for three cases: pure L, pure T and equal amount of L and T excitations.

An indirect possibility of performing polarization experiments is to look for decay asymmetry. Indeed the angular distribution of the decay pions from the resonance is a signature of the Δ polarisation. Such experiments have been recently undertaken at the Saturne laboratory with the Diogene detector. It is easy to show by a schematic calculation how one can determine which is the favoured L or T spin channel in a given experiment. We assume a driving interaction depending only on the direction of the momentum transfer $q$:

$$V_{eff} = F_L \sigma_1 \Delta \otimes \Delta + F_T (\sigma_1 \times \Delta) \otimes (\sigma_2 \times \Delta)$$

The angular distribution of the decay pions is then:

$$d\sigma/d\Omega \sim (F_L^2 + 2F_T^2) + (F_L^2 - F_T^2)P_3(\cos(\theta)) \sim 1 + [(F_L^2 - F_T^2)/(F_L^2 + 2F_T^2)]P_3(\cos(\theta))$$

(7)

As can be seen from Fig. 7 the distribution is concave or convex for dominant SL or ST components respectively and flat for equal coefficients (which means in conventional notations a ratio $W_{L/T} = 1/3$). More generally the interaction depends on the average projectile-ejectile momentum and the other polarization tensors $t_{21}$ and $t_{22}$ no longer vanish.

6 - CONCLUSION

There are all probabilities that the $(^3He, t)$ reaction has provided the first evidence of the nuclear pionic mode. This conclusion is reinforced by the decrease of the shift observed at higher projectile energy. This result insures that the models developed for the description of Δ's and π's in nuclei are valid at usual nuclear densities. However the problems raised by the $(d,2p)$ measurement prevent definitive commitment. In view of the crucial role of the $W_{L/T}$ balance, polarization experiments should be undertaken either with other polarized projectiles ($p$, $^3Li$,..) or with measurements of decay asymmetry. A number of experiments once proposed to measure the nuclear pionic field would also be new possible detectors of the pionic branch. Letting apart the challenging neutrino reactions one can quote a number of photopion reactions ($\pi, 2\gamma$), ($e, e'\pi^\pm$), ($\pi^\pm, e^+ e^-$) including the promising annihilation reaction/23-24/ of two pionic modes $\pi^+ + \pi^- \rightarrow e^+ + e^-$

The identification of the nuclear pionic mode is interesting by itself as a new "collective" mode of the nucleus at high excitation energy. It has far reaching consequences as giving a solid basis for extrapolation to higher nuclear densities as reached in heavy ions reactions. The pion yield and the stopping power of nuclear matter which are important quantities for the study of the equation of state are strongly affected by the softening of the pion dispersion line/25/ exemplified at normal density on Fig. 3. More tentatively following the evolution of the mode with ever increasing density may be informative on the way chiral symmetry is restored.
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