сообщения
объединенного
института
ядерных
исследований
дубна

E4-83-543

R.A.Eramzhyan, M.Gmitro, T.D.Kaipov,
S.S.Kamalov, R.Mach

INELASTIC ELECTRON SCATTERING AND RADIATIVE PION CAPTURE TO THE LOWEST 1^+ AND 2^+ ISOVECTOR LEVELS IN A=12 NUCLEI. CONTINUITY–EQUATION EFFECTS

1983

1. INTRODUCTION

Investigations of the pion photoproduction, pion inelastic scattering and radiative pion capture on light nuclei performed recently have led to a rather successful picture for all these processes taken separately. Obviously, one should try now a more ambitious task, namely try to understand these (and maybe other as well) reactions simultaneously taking as an input fixed nuclear structure information (transition densities) which provide an accurate description of the electron scattering data. There exist already a few examples where experimental results are available for (e, e′), (π, π′), (γ, π), at rest (π^-, γ) and (μ^-, ν_μ) transitions between the same pairs of levels.

In the present paper we shall study the isovector M1 and E2 transitions in A = 12 system, which connect the $0^+ T=0$ ground state of ^{12}C with 1^+ T = 1 (15.1 MeV) and 2^+ T = 1 (16.1 MeV) levels in ^{12}C or their isobar analogues (ground and 1st excited states) in ^{12}B and ^{12}N. These transitions have become an important testing ground for numerous prominent problems like the structure of the weak interactions [1-3], critical "nuclear opalescence" [4], and in the realm of nuclear physics itself for the problems of mesonic exchange currents [5]. Clearly, a safer understanding of the nuclear structure features of such widely utilised examples is highly desirable. Our objective is to extract the nuclear transition densities for the two levels from data on (e,e′) transitions and to examine to what extent the mechanisms of the pionic reactions proposed recently [6,7] are valid in different energy regions.

Excellent electron-scattering data for the M1 transition over a broad range of Q [5] and for the E2 transition for $Q \lesssim 0.5$ fm^{-1} (ref. [8]) and $0.5 < Q < 2.2$ fm^{-1} (refs. [9,10]) in ^{12}C have appeared a few years ago. They not only cover with a high precision a wide interval of the momentum transfer Q, but also provide (for E2) a separation of the longitudinal and transverse form factors. The interpretation of the data met, however, considerable difficulties.

The main problem is indeed the separation of spin-magnetisation (operator $e\vec{\mu}(\vec{r})$) and convection-current (operator $e\vec{j}(\vec{r})$) contributions to the transversal form factors $F_T(J)$. The 1p-shell configuration-mixing models (e.g. the Cohen-Kurath model used in ref. [9]) attribute all the E2-transverse strength to the spin operator $\vec{\mu}(\vec{r})$ simply because the convection-current contribution to $F_T(2^+)$ vanishes due to the identical geometry of the harmonic oscillator radial wave functions in the 1p-shell. The nuclear structure information extracted in this way, if used as an input in the further calculation, overestimates strongly, e.g., the pion photoproduction cross section for the $2^+(T=1,\ 0.95$ MeV) level in ^{12}B, see ref./11/. Similarly, the yield of the hard gamma rays from the radiative pion capture in A = 12 nuclei is always/12/ predicted incorrectly for the two lowest levels in ^{12}B: $[R(2^+)/R(1^+)]_{theor.} = 2[R(2^+)/R(1^+)]_{exper.}$ Also, inspecting the results for $|F_T(2^+)|^2$ at higher Q (ref./9/) together with the precision measurement of $|F_L(2^+)|^2$ and $|F_T(2^+)|^2$ near threshold/8/, one observes that the 1p-shell configuration-mixing models cannot account for both the sets of data simultaneously.

We argue that the major omission in the earlier work is connected with the form of the impulse approximation employed. Though it is known that using the continuity equation for the electromagnetic current one obtains a much more realistic form of the elementary amplitude to be used in the impulse approximation, this practice has been followed in the long wavelength regime only. Apparently, nobody has investigated yet the constraint due to the continuity equation in the electric transitions in ^{12}C. Numerically the effect is enormous. We shall see in sect.2 that after applying the continuity-equation constraint (CEC) for the calculation of $F_T(2^+)$ one can quantitatively reproduce the $|F_T(2^+)|^2$ and $|F_T(2^+)|^2$ form factors over the whole observed range of Q using the Cohen-Kurath wave function with a simple renormalization.

Simultaneously, the corresponding transition densities account correctly for the capture rates of the radiative pion capture from the mesoatomic orbitals. The results are displayed in Sect.3. The (π,π') and $(\gamma,\pi^\pm)$ reactions leading to the excitations of the lowest 1^+ (T = 1) and 2^+ (T = 1) levels in A = 12 nuclei will be considered in the subsequent publications.

2. NUCLEAR TRANSITION DENSITIES

The electron excitation of the 1^+ and 2^+ levels in nucleus with zero-spin ground state is fully characterised by three form factors: $F_T(1^+)$, $F_L(2^+)$, and $F_T(2^+)$. In the plane-wave Born approximation the form factors squared at momentum transfer Q are/13/

$$|F_L(J)|^2 = \frac{4\pi}{Z^2}\,\frac{|\langle J_n\,||\,j_J(Qr)\,Y_J(\Omega_{\vec r})\,||\,J_0\rangle|^2}{2J_0+1}\,f_p^2 f_{c.m.}^2 \qquad (1a)$$

$$|F_T(J)|^2 = \frac{4\pi}{Z^2}\,\frac{|\langle J_n\,||\,U_J(Q)+V_J(Q)\,||\,J_0\rangle|^2}{2J_0+1}\,f_p^2 f_{c.m.}^2 \qquad (1b)$$

with

$$U_{JM}(Q,el) = \frac{1}{Q}\int d^3r\,\vec\nabla\times[j_J(Qr)\,\vec Y^M_{JJ}(\Omega_{\vec r})]\cdot\hat{\vec j}_N(\vec r), \qquad (2a)$$

$$V_{JM}(Q,el) = Q\int d^3r\; j_J(Qr)\,\vec Y^M_{JJ}(\Omega_{\vec r})\cdot\hat{\vec\mu}_N(\vec r) \qquad (2b)$$

for the electric excitations, and

$$U_{JM}(Q,magn) = \int d^3r\; j_J(Qr)\,\vec Y^M_{JJ}(\Omega_{\vec r})\cdot\hat{\vec j}_N(\vec r) \qquad (3a)$$

$$V_{JM}(Q,magn) = \int d^3r\;\vec\nabla\times[j_J(Qr)\,\vec Y^M_{JJ}(\Omega_{\vec r})]\cdot\hat{\vec\mu}_N(\vec r) \qquad (3b)$$

for the magnetic ones. The total nuclear current density is

$$\hat{\vec J}(\vec r) = \hat{\vec j}_N(\vec r) + \vec\nabla\times\hat{\vec\mu}_N(\vec r), \qquad (4)$$

all unexplained notations are that of deForest and Walecka/13/. The single nucleon correction f_p^2 and centre-of-mass factor $f_{c.m.}^2$ are taken into account/5/.

Let us assume that the configuration mixing involves the $0p_{3/2}$ and $0p_{1/2}$ subshells only. Four coefficients a^J_{ij}, i, j = 1/2, 3/2 (see eq. (5) below) universal for the transition of the given tensorial rank J contain then the full nuclear structure information. The most popular model of this type is indeed the Cohen-Kurath model (CK)/14/. The model has generally been successful in describing the energy levels and transition probabilities of the p-shell nuclei. Here we face, however, a non-trivial extension to the processes with a rather high momentum transfer. In the rest of this section we shall discuss the validity and necessary modification of the model in our case. The coefficients a^J_{ij} which we use are displayed in table 1.

The M1 form factor in ^{12}C has received a very detailed attention in the last few years/3,5,15/. By now it is established that CK wave functions within impulse approximation (IA) fail to reproduce the M1 data in the interval Q > 1 fm^{-1}. Our calculations of the (π,π') and $(\gamma,\pi^\pm)$ differential cross sections have shown that this inadequacy of input indeed deteriorates the corresponding results at higher energies and large pion angles, where the typical momentum transfer amounts to $Q \approx 1.5\div2$ fm^{-1}.

Table 1

Values of harmonic oscillator parameter b (fm) and the coefficients a^J_{ij} used to parametrize the nuclear transition densities. The Cohen-Kurath wave functions correspond to their (8-16) 2BME interaction

Transition	Set	b	$a_{\frac{3}{2}\frac{3}{2}}$	$a_{\frac{3}{2}\frac{1}{2}}$	$a_{\frac{1}{2}\frac{3}{2}}$	$a_{\frac{1}{2}\frac{1}{2}}$	Ref.
$0^+ \to 1^+$	CK	1.888	−0.083	−0.327	−0.691	−0.073	11,14
$0^+ \to 1^+$	DH	1.757	−0.086	0.0	−0.338	0.111	5,11
$0^+ \to 2^+$	CKS[a)]	1.64	−0.041	0.076	−0.456	0.0	11,14

a) Coefficients $a^{J=2}_{ij}$ given here are scaled down by a factor of $\sqrt{0.45}$ as compared with those of refs./11,14/, see the text.

An interesting development is due to Dubach and Haxton /5/ (DH). They have shown that the phenomenological form $|F_T(1^+)|^2 \propto Q^2 \exp(-2y)(A+By) f_p^2 f^2_{c.m.}$, where $y = (Qb/2)^2$ and b is the oscillator parameter, provides a very good fit to the M1 form factor over the whole observed range of Q. We shall use their phenomenological coefficients as well.

The case of E2 transition is much more complicated. The Cohen-Kurath wave function requires a very strong downward renormalisation in order to account for the experimental B(C2) value extracted from the longitudinal form factor $|F_L|^2$ in ref./8/. Friebel et al./8/ ascribe this feature to the omission of the higher ($2\hbar\omega$, etc.) nuclear configurations in the Cohen-Kurath model. They have performed several configuration-mixing calculations which included some selected $2\hbar\omega$ configurations and reduced considerably the theoretical B(C2) value. They have also suggested that the effect can approximately be viewed as a scaling

$$\langle J_n, T_n T_{zn} \| \mathcal{O}_J \mathcal{P}_{TT_z} \| 0,00 \rangle = \frac{\beta}{\sqrt{2T_n+1}} \sum_{i,j=\frac{1}{2},\frac{3}{2}} a^J_{ij} \langle p_i \||\mathcal{O}_J \mathcal{P}_T\|| p_j \rangle \qquad (5)$$

and have estimated via perturbation theory the value $\beta^2 = 0.40$. (The s.p. matrix element $\langle p_i \||\mathcal{O}_J \mathcal{P}_T\|| p_j\rangle$ is reduced both in the usual and isospin space). In contradiction with this a strong upward renormalization ($\beta^2 \simeq 2$) would be needed to reproduce the transversal form factor $|F_T(2^+)|^2$ at $Q < 0.5$ fm^{-1} reported in ref./8/. At the same time Flanz et al./9/ with the CK

wave functions have reproduced the form factor $|F_T(2^+)|^2$ in the domain of a higher momentum transfer.

We think that the contradictions here are due to a straightforward use of eqs. (2). As a matter of fact, for a more reliable calculation of the transverse electric form factor one should apply the continuity equation

$$\vec\nabla\cdot\hat{\vec{j}}_N(\vec r) = -\frac{\partial\hat\rho_N}{\partial t} = -i\,[\hat H, \hat\rho_N(\vec r)] \qquad (6)$$

for the nuclear current density $\hat{\vec{j}}_N$ and the nuclear change density $\hat\rho_N$. For the sake of comparison we first quote the standard result. After the usual manipulations /13/ the impulse-approximation matrix element of eq. (2a) acquires the form

$$\langle J_n \| U_J(Q,el) \| J_0 \rangle = \sum_{j=1}^{A} \frac{e_j}{M} \langle J_n \| -\left(\frac{J}{2J+1}\right)^{1/2} j_{J+1}(Qr_j)\vec Y_{JJ+1}(\Omega_{\vec r_j})\cdot\vec\nabla + \qquad (7)$$

$$+\left(\frac{J+1}{2J+1}\right)^{1/2} j_{J-1}(Qr_j)\vec Y_{JJ-1}(\Omega_{\vec r_j})\cdot\vec\nabla \| J_0 \rangle .$$

To apply CEC we first rewrite eq. (2a) in terms of vector harmonic $\vec Y^M_{JJ-1}$ and $\vec Y^M_{JJ+1}$ as

$$i\int d^3r \left[\left(\frac{J+1}{2J+1}\right)^{1/2} j_{J-1}(Qr)\vec Y^M_{JJ-1}(\Omega_{\vec r_j}) - \left(\frac{J}{2J+1}\right)^{1/2} j_{J+1}(Qr)\vec Y^M_{JJ+1}(\Omega_{\vec r})\right]\cdot\hat{\vec j}_N(\vec r) \qquad (2a')$$

and then, using the identity

$$j_{J-1}(Qr)\vec Y^M_{JJ-1} = \frac{1}{Q}\left(\frac{2J+1}{J}\right)^{1/2}\vec\nabla\, j_J(Qr)\, Y_{JM}(\Omega_{\vec r}) - \left(\frac{J+1}{J}\right)^{1/2} j_{J+1}(Qr)\vec Y^M_{JJ+1}(\Omega_{\vec r})$$

eliminate $\vec Y^M_{JJ-1}$. Integrating by parts the term with the gradient operator and using the continuity equation (6) we find the final result for the IA-matrix element:

$$\langle J_n \| U^{CEC}_J(Q,el) \| J_0 \rangle =$$

$$= \sum_{j=1}^{A} e_j \langle J_n \| -\left(\frac{J+1}{J}\right)^{1/2} \frac{E_n - E_0}{Q} j_J(Qr_j)\, Y_J(\Omega_{\vec r_j}) - \qquad (8)$$

$$- \frac{1}{M}\left(\frac{2J+1}{J}\right)^{1/2} j_{J+1}(Qr_j)\, \vec Y_{JJ+1}(\Omega_{\vec r_j})\cdot\vec\nabla \| J_0 \rangle .$$

For details we refer to the recent review by Heisenberg and Blok /16/.

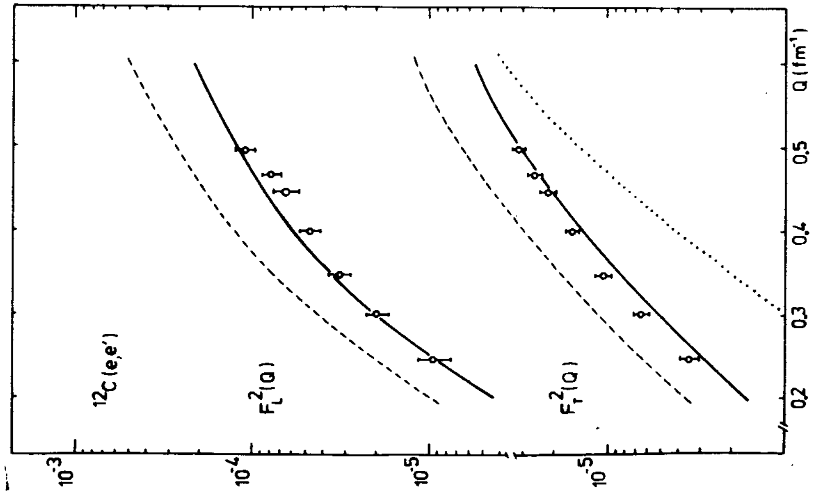

Fig.1. The inelastic form factors for the 16.1 MeV 2^+ T = 1 level in ^{12}C. The dotted line corresponds to the calculation using eq. (7), i.e., without CEC. All curves are obtained with the CK wave function either unmodified (dashed line) or scaled as in table 1 (solid line). Data are from ref. /8/

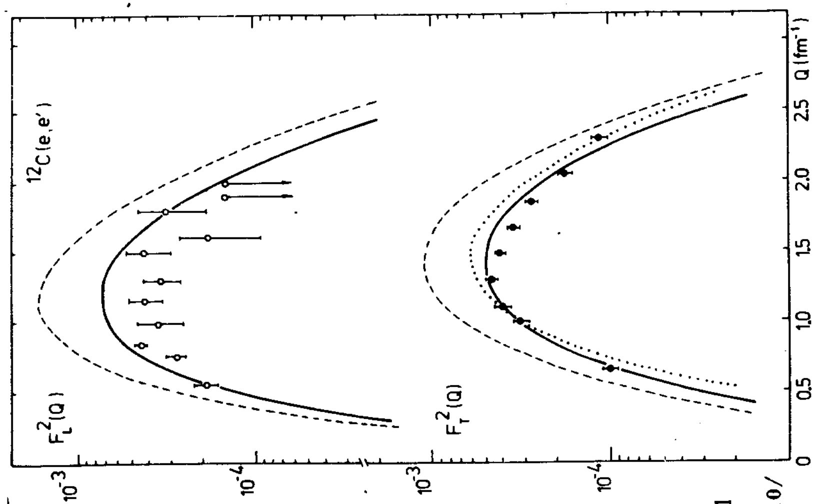

Fig.2. The same as in fig.1 but data are taken from ref. /9/ ($|F_T|^2$) and ref. /10/ ($|F_L|^2$).

Physically the procedure is connected with the fact that the impulse-approximation form of the nuclear charge

$$\hat{\rho}_N(\vec r) = \sum_{j=1}^{A} e_j \, \delta(\vec r_j - \vec r) \qquad (9)$$

is expected to be more reliable than the corresponding expression

$$\hat{\vec j}_N(\vec r) = \sum_{j=1}^{A} [\delta(\vec r_j - \vec r) \frac{e_j}{iM} \vec\nabla\,]_{sym.} \qquad (10)$$

for the nuclear current. In other words, the form (8) includes implicitly a large portion of the effects connected with the mesonic exchange currents, which otherwise should have been explicitly calculated for the form (7). The argumentation has been suggested already at the beginning of nuclear physics in connection with the well-known Siegert theorem /17/. Note, however, that we do not use the long wavelength (LWL) regime as it was done, e.g., by Ueberall and collaborators in their calculation with the Helm model /18/. The LWL form can be obtained from eq. (8) by omitting the second term in the matrix element which is indeed not small except maybe for $Q < 0.1$ fm^{-1}.

The experimental and calculated form factors for the 2^+ (T = 1, 16.1 MeV) level are shown in figs. 1 and 2. It is easily seen that using the continuity-equation constraint displayed by eq.(8) to calculate $|F_T(2^+)|^2$ one can bring both $|F_L(2^+)|^2$ and $|F_T(2^+)|^2$ to a very reasonable agreement with data for a broad interval of Q just by scaling the Cohen-Kurath wave function by a factor $\beta = \sqrt{0.45}$, see the dashed and full lines in figs.1 and 2.

A word of warning is needed at this place. Certainly, the 2^+ T = 1 (16.1 MeV) level must be a bit lucky example. We do not expect that CEC alone can explain with a simple scaling parameter all electric transitions in the p-shell nuclei if a strongly limited ($0p_{3/2}$, $0p_{1/2}$ only) configuration mixing is prescribed. For example, the lowest excited state in ^{12}C (2^+ T = 0, 4.44 MeV) seems to be of a more complicated nature and besides CEC also the $2\hbar\omega$ components of the wave functions would be needed to account for the F_L and F_T form factors of this level as reported by Flanz et al. /9/.

The nuclear structure model like CK is indeed preferable even if it forces us to fit a parameter like β. The wave function still retains the information which allows to reconstruct

the individual contributions of $\hat{\vec j}_N(\vec r)$ and $\hat{\vec\mu}_N(\vec r)$ terms. This information is lost in the phenomenologically fitted transition densities. Unlike the case of M1 transition, for the E2 transition both convection-current and spin-magnetisation terms contribute appreciably to F_T. Their relative importance in the pionic processes considered below is indeed different and we shall see in sect.3 an example of the possible effects.

3. RADIATIVE PION CAPTURE

The yield of the hard gamma rays due to the capture of the negative pions resting on the mesoatomic orbitals is an important source of information about the nuclear spin densities. For a recent review of the (π^-,γ) reaction see ref./19/. The partial transition rates in many p-shell nuclei have been measured with a very good precision. As for our case, the repeated careful calculations/12/ of the partial transitions in A = 12 nuclear system have shown a standing discrepancy: The calculated yield for the 2^+ first excited state in ^{12}B comes out systematically too high, by a factor of at least two, though the 1^+ ground-state yield is obtained very close to its experimental value.

In this work we have calculated the capture rates $\Lambda(n_\pi \ell_\pi)$ for the radiative capture from the mesoatomic orbital $\phi_{n_\pi \ell_\pi}$ using the momentum-space representation. This helps much in performing the relativistic transformation/7/ between the pion-nucleon and pion-nucleus c.m. systems. Note that Nagl et al./20/ found this transformation to be numerically very important at the threshold of the pion photoproduction reaction (see fig.8 of ref. /20/). Using the partial-wave photoproduction amplitudes/7/ $U^{\gamma j}(q\ell_\pi, k\lambda\ell_\gamma)$ one obtains for the radiative capture rate

$$\Lambda(n_\pi \ell_\pi) = \frac{1}{(2J_0+1)(2\ell_\pi+1)} \; \frac{k}{2\pi^2 M_{\pi A}} \sum_{j\lambda\ell_\gamma} \frac{2j+1}{2\ell_\gamma+1} |B(n_\pi \ell_\pi)|^2 , \qquad (11)$$

where J_0 is the total spin of the nuclear initial state, the pion-nucleus reduced mass if $M_{\pi A} = m_\pi M_A/(m_\pi + M_A)$, and

$$B(n_\pi \ell_\pi) = \int \tilde{U}^{\gamma j}_{n0}(k\lambda\ell_\gamma, q\ell_\pi)\, \phi_{n_\pi \ell_\pi}(q)\, q^2 dq \qquad (12)$$

with k and λ for the photon momentum and polarization. Here we have introduced the notation $\tilde{U}^{\gamma j}_{n0}$ for the time-reversed photoproduction amplitude connected with the transition to the nuclear state $|n\rangle$. Pionic wave functions in the momentum space are taken as

$$\phi_{n_\pi \ell_\pi}(q) = \sqrt{C_{n_\pi \ell_\pi}}\, \phi^{Coul.}_{n_\pi \ell_\pi}(q), \qquad (13)$$

where $\phi^{Coul.}_{n_\pi \ell_\pi}$ is the Coulomb wave function for a point-charge, and the coefficients $C_{n_\pi \ell_\pi}$ take care of the finite size and strong interaction effects. It has been shown/19/, that with $C_{1s} = 0.45$ and $C_{2p} = 1.4$ one obtains very good approximation of the functions $\phi_{n_\pi \ell_\pi}$.

The measured radiative branching ratio appears as the sum of partial branching ratios for the different orbits:

$$R = \omega_s \frac{\Lambda(1s)}{\Gamma(1s)} + \omega_p \frac{\Lambda(2p)}{\Gamma(2p)} , \qquad (14)$$

where the mesoatomic absorption strength ω_{ℓ_π} and the total absorption rate $\Gamma(n_\pi \ell_\pi)/\hbar$ are taken at their experimental values/19/ as $\omega_s = 0.08$, $\omega_p = 0.92$, $\Gamma(1s) = 3.12 \times 10^{-3}$ MeV, and $\Gamma(2p) = 1.25 \times 10^{-6}$ MeV.

Table 2

Comparison of measured partial (π^-,γ) branches with the shell model predictions for A = 12

Transition	$R^{a)}_{exp}$ (10^{-4})	Set$^{b)}$	R_{theor} (10^{-4})
$0^+ \to 1^+$	6.22 ± 0.35	CK	6.32
$0^+ \to 2^+$	1.29 ± 0.25	CKS	1.57

a) Ref./21/.

b) See table 1.

In table 2 we present our results for $R(1^+)$ and $R(2^+)$ calculated with the transition densities as parametrized in table 1, and the available data. Using the modified CK wave function of the 2^+ level we have obtained for the first time a decent agreement of calculated and experimental yields R. This is indeed due to the use of CEC in the analysis of (e,e') data. The (unmodified) CK wave function of the 2^+ level, if used in eq. (7), would reproduce correctly the value of $|F_T(2^+)|_{Q^2 \approx m_\pi^2}|^2$, but it ascribes the transition totally to the spin-magnetisation operator. This would then lead to an overestimate of the $R(2^+)$ value as encountered in the previous works. Via eq. (8) important convection-current contributions appear in the E2 transition thus reducing the role of the spin-magnetization term with the needed effect on the calculated $R(2^+)$ value.

REFERENCES

1. Mukhopadhyay N.C. Phys.Rep., 1977, 30C, p.1.
2. Balashov V.V., Korenman R.A., Eramzhyan R.A. Meson Capture by Atomic Nuclei. Atomizdat, M., 1978.
3. Roesch L.Ph. et al. Phys.Rev.Lett., 1981, 46, p.1507; Phys.Lett., 1981, 107B, p.31.
4. Delorme J. et al. Phys.Lett., 1980, 89B, p.327; Delorme J., Figureau A., Guichon P. Phys.Lett., 1981, 99B, p.187.
5. Dubach J., Haxton W.C. Phys.Rev.Lett., 1978, 41, p.1453; Flanz J.B. et al. Phys.Rev.Lett., 1979, 43, p.1922; Haxton W.C. Phys.Lett., 1978, B76, p.165.
6. Gmitro M., Mach R. Zeit.f.Phys., 1979, A290, p.179; Gmitro M., Kvasil J., Mach R. Phys.Lett., 1982, 113B, p.205.
7. Eramzhyan R.A. et al. J.Phys.G: Nucl.Phys., 1983, 9, p.605.
8. Friebel A. et al. Nucl.Phys., 1978, A294, p.129.
9. Flanz J.B. et al. Phys.Rev.Lett., 1978, 41, p.1642.
10. Yamaguchi A. et al. Phys.Rev., 1971, C3, p.1750.
11. Singham M.K., Tabakin F. Ann.Phys., 1981, 135, p.71.
12. Kissener H.R. et al. Nucl.Phys., 1978, A302, p.523; Ohtsubo H., Nishiyama T., Kawaguchi M. Nucl.Phys., 1974, A224, p.164.
13. de Forest T. Jr., Walecka J.D. Adv.Phys., 1966, 15, p.1.
14. Cohen S., Kurath D. Nucl.Phys., 1965, 73, p.1; 1967, A101, p.1
15. Mukhopadhyay N.C., Martorell J. Nucl.Phys., 1978, A296, p.461; Suzuki T. et al. Phys.Lett., 1981, 106B, p.19.
16. Heisenberg J., Block H.P. Ann.Rev.Nucl.Part.Sci., 1983, 33, p.2224.
17. Sachs R.G. Nuclear Theory. Addison-Wesley, Cambridge, Mass., 1953.
18. Nagl A., Überall H. Phys. Lett., 1980, 96B, p.254.
19. Gmitro M. et al. Sov.J.Part. and Nuclei, 1982, 13, p.1230; 1983, 14, p.879.
20. Nagl A., Cannata F., Überall H. Acta Phys.Austr., 1978, 48, p.267.
21. Perrenoud A. et al. Report, Université de Lausanne, 1981.

Received by Publishing Department
on July 27, 1983.

SUBJECT CATEGORIES OF THE JINR PUBLICATIONS

Index	Subject
1.	High energy experimental physics
2.	High energy theoretical physics
3.	Low energy experimental physics
4.	Low energy theoretical physics
5.	Mathematics
6.	Nuclear spectroscopy and radiochemistry
7.	Heavy ion physics
8.	Cryogenics
9.	Accelerators
10.	Automatization of data processing
11.	Computing mathematics and technique
12.	Chemistry
13.	Experimental techniques and methods
14.	Solid state physics. Liquids
15.	Experimental physics of nuclear reactions at low energies
16.	Health physics. Shieldings
17.	Theory of condenced matter
18.	Applied researches
19.	Biophysics

Эрамжян Р.А. и др. E4-83-543

Неупругое рассеяние электронов и радиационный π-захват на низшие 1^+ и 2^+ изовекторные состояния ядер с $A = 12$. Эффекты, связанные с уравнением непрерывности

Исходя из уравнения непрерывности для электромагнитного тока ядра $\vec{\nabla}\cdot\vec{J}_N = -\partial\rho_N/\partial t$ проведен анализ переходных плотностей в ^{12}C. Полученные результаты значительно отличаются от результатов предшествующих анализов E2-перехода. Зафиксированы плотности M1 и E2 переходов, которые правильно описывают данные по электронному рассеянию. Полученная таким образом информация о ядерной структуре используется для расчета скоростей радиационного захвата пионов.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

Сообщение Объединенного института ядерных исследований. Дубна 1983

Eramzhyan R.A. et al. E4-83-543

Inelastic Electron Scattering and Radiative Pion Capture to the Lowest 1^+ and 2^+ Isovector Levels in $A = 12$ Nuclei. Continuity-Equation Effects

Continuity equation for the nuclear electric charge and convection current $\vec{\nabla}\cdot\vec{J}_N = -\partial\rho_N/\partial t$ has been used in an analysis of nuclear transition densities in ^{12}C. The results differ considerably from the former derivations. Standard M1 and newly calculated E2 nuclear transition densities are fixed which provide an accurate description of the electron scattering data. Such a nuclear structure imput is used in the radiative pion capture calculations.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

Communication of the Joint Institute for Nuclear Research. Dubna 1983