28/IV-84

сообщения
объединенного
института
ядерных
исследований
дубна

E4-84-37

M.Kirchbach, S.Kamalov, H.-U.Jäger*

SHORT-RANGE AND MEDIUM POLARIZATION EFFECTS IN THE TWO-BODY WEAK AXIAL CHARGE DENSITY FOR THE ISOVECTOR TRANSITIONS $^{16}O(0^+_1) \leftrightarrow {}^{16}N(0^-_1)$

* Central Institute for Nuclear Research Rossendorf, 8051 Dresden, PO Box 19, GDR

1984

1. INTRODUCTION

The β-decay and μ^--capture experiments in the A = 16 nuclear system were recently studied in a series of papers [1-3] as one of the most confident sources of getting information on the relevance of the meson exchange contribution to the time component J^4 of the nuclear weak axial vector current. The longstanding discrepancy between the experimental value of the ratio of the partial muon capture rate $\Lambda^{\mu}(0_1^+ \to 0_1^-)$ to the partial β-decay rate $\Lambda^{\beta}(0_1^- \to 0_1^+)$ and its theoretical prediction based on the one-body current $J^w_{(0)}$ was acceptably explained in terms of an enhancement of the one-body weak-axial charge density due to the long-range one-pion exchange current $J^4_{(\pi)}$. In our previous paper [3] we have exploited the phenomenological Lagrangian version of the hard pion model [4] to show that a consistent description of the virtual meson exchange in the nucleus leads to a non-negligible contribution from short-range heavy (rho- and A_1) meson exchanges. The latter have been shown to damp the nuclear matrix element (n.m.e.) of the exchange operator. However, to get a consistent description of the nuclear structure and the mesonic exchange current (MEC) contribution to the nuclear observables of interest, we have to complete the short-range effects by calculating also corrections due to repulsive two-body correlations in the wave functions of the initial and final nuclear states. The extended spatial structure of the meson-nucleon vertices of the exchange current combined with the effect of the medium polarization are also considered. In addition, the sensitivity of the nuclear matrix element of the current operator to the oscillator strength parameter b is studied in order to shed some light on one source of uncertainty in the comparison of the theoretical to the experimental [5,6] partial β-decay rate.

2. THE TWO-BODY WEAK AXIAL CHARGE DENSITY OPERATOR

The two-body weak axial charge density operator used (it will be referred to as $J^4_{(\pi')}$) has been constructed in ref. [3] in the one-boson approximation by means of the S-matrix method and contains, except for a long-range one-pion exchange part also a piece (denoted by $J^4_{(\rho A_1)}$) which is due to short-range ρ- and A_1-meson exchanges. We also include the contribution from the $\Delta(1236)$

isobar excitation current restricting ourselves to the terms $\delta_{\lambda\mu} - 1/3\,\gamma_\lambda\gamma_\mu$ in the propagator. We completely neglect the momentum dependent terms as well as the $\rho N\Delta$-exchange diagram. A discussion on the validity of this approximations can be found elsewhere, for example, in ref. /7/. In the hard pion model the MEC operator reads

$$J^4_{(\pi')}(1,2) = \frac{g_A g_\rho^2}{2m_\rho^2}\left(1+\frac{4f^2_{\pi NN^*} f_\pi^2}{9M(M^*-M)m_\pi^2}\vec q^{\,2}\right)\frac{\vec\sigma_2\cdot\vec q}{m_\pi^2+q^2}(\vec\tau_1\times\vec\tau_2)_\mp +$$

$$+ \frac{g_A g_\rho^2}{2}\left(1-\frac{\kappa_V}{4M^2}\vec q^{\,2}\right)\left(1-\frac{1}{m_\rho^2}\vec q^{\,2}\right)\frac{\vec\sigma_2\cdot\vec q}{((q+k)^2+m_\rho^2)(m_\pi^2+q^2)}(\vec\tau_1\times\vec\tau_2)_\mp + (1\leftrightarrow 2). \qquad /1/$$

Here $f_{\pi NN^*}$ stands for the πNN^* coupling constant with $f^2_{\pi NN^*}/4\pi = 0.35$, f_π denotes the pion decay constant ($f_\pi = 92$ MeV), g_ρ is the ρ-meson coupling constant ($2f_\pi^2 g_\rho^2 = m_\rho^2$). The Δ-isobar, the nucleon, the pion and the ρ-meson masses are M^*, M, m_π and m_ρ, respectively and κ_V is the tensor coupling constant of the ρ-meson ($\kappa_V = 3.7$). The pion decay constant is connected with the πNN-coupling constant g_r ($g_r^2/4\pi = 14.6$) by the Goldberger-Treiman relation $g_r f_\pi = Mg_A$, where g_A is the axial form factor of the nucleon. The term with the denominator $1/M(M^*-M)$ represents a small contribution from the Δ-isobar excitation current and will enhance the leading one-pion exchange term by 1.5%. The Fourier transform of the one-meson exchange term $\vec\sigma\cdot\vec q/(m^2+q^2)$ will result in $im^2Y_1(mr)$, were the Yukawa function is defined as $Y_1(x) = (1+1/x)\exp(-x)/x$, $\vec k$ and $\vec q$ are the respective current and pion momenta.

To include the renormalization effects due to the finite range of the meson nucleon vertices and to the medium polarization we completely adopt the treatment of Riska et al. /8,9/, where the meson propagators are modified by an additional self-evergy to the form $(m^2_{(\pi,\rho)}+q^2+\Pi^{(\pi,\rho)})^{-1}$. Further, the function $U^{(\pi,\rho)}(q)$ is introduced by the following relation to the self-energy $\Pi_0^{(\pi,\rho)}$ of static mesons, $\Pi_0^{(\pi,\rho)} = -q^2U^{(\pi,\rho)}(q)$. For a simple form $R = g'\delta_{mn}$ of the pion-irreducible isobar-hole interaction the self-energy is written as $\Pi(q) = \Pi_0(q)/(1+g'U(q))$, where a value $g' = 0.6$ is adopted. In addition, at each vertex πNN, $\pi N\Delta$, and $\pi N\rho$ of the diagram representation of the meson exchange current (see the figure in ref. /8/) medium renormalized monopole hadronic form factors are introduced,

$$F_{(\pi,\rho)}(q) = \frac{\Lambda^2_{(\pi,\rho)} - m^2_{(\pi,\rho)}}{(\Lambda^2_{(\pi,\rho)}+q^2)(1+g'U^{(\pi,\rho)}(q))}. \qquad /2/$$

For the mass scale parameters the following values are used, $\Lambda_\pi = 1.18$ GeV/c^2 and $\Lambda_\rho = 2.0$ GeV/c^2. At the lepton-hadron vertices the bare hadronic form factors are to be incerted. Finally, nuclear matter calculations lead to $\Pi_0^\rho(q) \approx -1.1\,q^2$ and $1 + g'U^\pi(q) \approx 1.53$. Within this scheme and in the soft-current limit $k \to 0$ the MEC operator in eq. /1/ becomes

$$J^4_{(\pi')}(1,2) = \frac{(\Lambda_\pi^2 - m_\pi^2)^2}{(\Lambda_\pi^2 - m_\pi^{*2})^2(1-(1-g')U^\pi(0))}\,\frac{g_A g_\rho^2}{m_\rho^2}\times$$

$$\times\left[\frac12 + \frac{m_\rho^2}{2(m_\rho^2-m_\pi^2)}\left(1-\frac{\kappa_V}{4M^2}\vec q^{\,2}\right)\left(1-\frac{1}{m_\rho^2}\vec q^{\,2}\right) + \frac{2f^2_{\pi NN^*}f_\pi^2}{9m_\pi^2 M(M^*-M)}\vec q^{\,2}\right]\vec\sigma_2\cdot\vec q\,\times$$

$$\times\left[\frac{1}{m_\pi^{*2}+q^2} - \frac{1}{\Lambda_\pi^2+q^2} - \frac{\Lambda_\pi^2-m_\pi^{*2}}{(\Lambda_\pi^2+q^2)^2}\right](\vec\tau_1\times\vec\tau_2)_\mp \; - \qquad /3/$$

$$- \frac{g_A g_\rho^2}{2(m_\rho^2-m_\pi^2)}\left(1-\frac{\kappa_V}{4M^2}\vec q^{\,2}\right)\left(1-\frac{1}{m_\rho^2}\vec q^{\,2}\right)\frac{(\Lambda_\rho^2-m_\rho^2)^2}{m_\rho^2(1+g'U^\rho(0))}\,\frac{\vec\sigma_2\cdot\vec q}{(\Lambda_\rho^2+q^2)^2}(\vec\tau_1\times\vec\tau_2)_\mp + (1\leftrightarrow 2),$$

$$m_\pi^{*2} = \frac{1+g'U^\pi(0)}{(1-(1-g')U^\pi(0))}\,m_\pi^2 .$$

After performing the transformation to the coordinate space we obtain

$$J^4_{(\pi')} = \frac12\sum_{i<j} e^{-i\frac{\vec k\cdot\vec R A}{2}} \sum_{+,-} \pm\left(e^{i\frac{\vec k\cdot\vec r A}{2}} \pm e^{-i\frac{\vec k\cdot\vec r A}{2}}\right)(\vec\sigma_i\pm\vec\sigma_j)\cdot\hat r\,\tilde\phi(r)\, i(\vec\tau_i\times\vec\tau_j)_\mp ,$$

$$A \equiv b\sqrt2,\quad \vec R = \frac{1}{b\sqrt2}(\vec r_i+\vec r_j),\quad \vec r = \frac{1}{b\sqrt2}(\vec r_i - \vec r_j), \qquad /4/$$

where b denotes the oscillator strength parameter in fm and the radial dependence is

$$\tilde\phi(r) = -\frac{(\Lambda_\pi^2-m_\pi^2)^2}{(\Lambda_\pi^2-m_\pi^{*2})^2(1-(1-g')U^\pi(0))}\,\frac{m_\pi^{*2}g_r^2}{8\pi M^2 g_A}\Big[\phi(m_\pi^{*2})Y_1(m_\pi^* r) - \qquad /5/$$

$$- \frac{\Lambda_\pi^2}{m_\pi^{*2}}\phi(\Lambda_\pi^2)Y_1(\Lambda_\pi r) - \frac{\Lambda_\pi^2-m_\pi^{*2}}{2m_\pi^{*2}}\phi(\Lambda_\pi^2)e^{-\Lambda_\pi r}\Big] + \frac{(\Lambda_\rho^2-m_\rho^2)^2 m_\rho^2 g_r^2}{2m_\rho^4(1+g'U^\rho(0))16\pi M^2 g_A}\times$$

$$\times\, c(\Lambda_\rho^2)e^{-\Lambda_\rho r},$$

$$\phi(y^2) = \frac{1}{2} + \frac{1}{2}\, c(y^2) + c_\Delta(y^2), \quad c(y^2)\ \frac{m_\rho^2}{m_\rho^2 - m_\pi^2}\left(1 + \frac{\kappa_V y^2}{4M^2}\right)\left(1 + \frac{y^2}{m_\rho^2}\right),$$

$$c_\Delta(y^2) = \frac{2 f^2_{\pi NN^*} f_\pi^2 y^2}{9M(M^* - M)\, m_\pi^2}.$$

3. EXPANSION OF THE NUCLEAR MATRIX ELEMENT OF THE CURRENT OPERATOR

The partial transition rates are proportional to the squared matrix element of the weak axial current J^w between the initial and final nuclear states

$$\Lambda^{(\mu,\beta)}(0_1^+ \rightleftarrows 0_1^-) \sim |\langle 0_1^-, T = 1 | J^w | 0_1^+, T = 0 \rangle|^2,$$

$$J^w = \sum_{i=1}^{A} J^w_{(0)}(i) + \sum_{i<j} J^4_{(\pi')}(i,j). \qquad /6/$$

Detailed expressions for $\Lambda^{(\mu,\beta)}$ and the respective one-body currents can be found in refs. /1-3/. In the second-quantization formalism the matrix element of a spherical tensor component $J^{(\lambda;t)}$ of the current operator in eq./4/ (λ and t are the respective multipolarities in the spin-coordinate and strong isospin spaces) reads

$$(f|||J^{(\lambda,t)}|||i) = \sum_{a,b} \hat{\lambda}^{-1}\, \hat{t}^{-1}\, (a|||J^{(\lambda,t)}_{(0)}|||b)\, \rho^{i\to f}_{[ab]\lambda,t} +$$

$$+ \sum_{\substack{a<b\\ c>d\\ J,T,J',T'}} \hat{\lambda}^{-1}\, \hat{t}^{-1}\, (ab;JT|||J^{(\lambda,t)}_{(\pi')}|||cd;J'T')\, \rho^{i\to f}_{[[ab]JT;[cd]J'T']\lambda,t}. \qquad /7/$$

The Roman alphabet indices denote a complete set of quantum numbers defining a single-particle state except for the magnetic quantum numbers, the brackets denote irreducible tensor product and $\hat{t} = (2t+1)^{1/2}$. The one- and two-body density matrices are defined similar to ref. /10/ as

$$\rho^{i\to f}_{[ab]\lambda,t} = (f|||[a_a^+ a_b]^{\lambda,t}|||i),$$

$$\rho^{i\to f}_{[[ab]JT;[cd]J'T']\lambda,t} = n_{ab} n_{cd}\, (f|||[[a_a^+ a_b^+]^{JT}\, [a_c a_d]^{J'T'}]^{\lambda,t}|||i), \qquad /8/$$

$$n_{\ell m} = (1 + \delta_{\ell m})^{1/2}.$$

In the present calculation they are generated by shell-model wave functions with configuration mixing which were obtained

on a harmonic oscillator single-particle basis by diagonalization of a nuclear residual interaction of Tabakin's type /11/. The ground state of ^{16}O consists of the closed core configuration (0p - 0h) and all $((2s1d)^1 1s^{-1})$, $((2p1f)^1 1p^{-1})$, $((2s1d)^2\ 1p^{-2})$ and $(2s^2 1d^{-2})$ configurations. The negative parity state contains up to 95% the particle-hole state $(2s^1 1p^{-1})$ plus seventy 2p-2h-components. In calculating the two-body density matrices it is sufficient to keep the two strongest (~1% each) configurations $((1d1f)^2\ 1p^{-2})$ and $(2s^2\ 1p^{-2})$ from the $3\hbar\omega$-space.

The repulsive hard core of the nucleon-nucleon interaction leads to strong two-body correlations which cannot be adequately reproduced by the shell-model configuration mixing in the truncated $2\hbar\omega(3\hbar\omega)$ space. That's why we will supplement the shell-model wave functions by multiplying them by a Jastrow-type cluster expansion restricting ourselves to two-body clusters, i.e., by $(1 - \sum_{i<j} f_{ij})$ /12,13/. We use in our calculation the parametrization of Miller and Spencer /14/

$$f_{ij} = f(r = |\vec{r}_i - \vec{r}_j|) = \exp(-\alpha r)(1 - \beta r^2),$$

$$\alpha = 1{,}1\ \mathrm{fm}^{-2}, \quad \beta = 0.68\ \mathrm{fm}^{-2}, \qquad /9/$$

which was checked in calculating MEC-like parity violating matrix elements /15,16/. According to the familiar Talmi-Moshinsky transformation of two-body matrix elements from jj-coupling scheme to the relative and centre-of-mass coordinates r and R of a pair of particles, the doubly reduced matrix elements in eq./6/ are determined by a linear combination (l.c.) of radial integrals of the type $(ab; JT|||J^4_{(\pi')}|||cd; J'T') \sim$ l.c. $\int_0^\infty dr\, r^2\, R_{n\ell}(r)\, j_L(kr/\sqrt{2})\, Y_1(m_{(\pi,\rho)} r)\, R_{n'\ell'}(r)$. Here $R_{n\ell}$ is the normalized spatial part of the harmonic oscillator wave function and j_L is the spherical Bessel function. The Jastrow correlations are introduced by the replacement $R_{n\ell}(r) \to (1 - f(r)) R_{n\ell}(r)$. Their influence on the n.m.e. of the one-body current can be treated in a similar way by exploiting the relation

$$\sum_{i=1}^{A} J^w_{(0)}(i) = \frac{1}{2(A-1)} \sum_{i<j} (1 - f(r))\,(J^w_{(0)}(i) + J^w_{(0)}(j))\,(1 - f(r)), \qquad /10/$$

with

$$J^w_{(0)}(i) = g^{(\mu,\beta)}\, \vec{J}_{(0)}(i) + J^4_{(0)}(i).$$

Here $\vec{J}_{(0)}(i)$ $(J^4_{(0)}(i))$ stands for the space-like (time-like) part of the one-body current. The constants g^μ and g^β are given by

$$g^{\mu} = g_A\sqrt{2}\,((g_p/g_A - 1)\frac{E_\nu}{2M} - 1),\quad g^{\beta} = g_A\, 0{,}932\sqrt{2}\,(1 + \frac{3\alpha Z}{2RE_0}),$$

$E_\nu = 95.121$ MeV, $E_0 = 11.05$ MeV, $R = 3.51$ fm, $\alpha = 1/137$.

Further,

$$\vec{J}_{(0)}(i) = j_1(qr_i)\vec{\sigma}_i\cdot\hat{r}_i\,\tau_-(i),\quad J^4_{(0)}(i) = \frac{\sqrt{2}\,g_A}{M} j_0(qr_i)\vec{\sigma}_i\cdot\vec{\nabla}_i\,\tau_-(i), \qquad /11/$$

$$\tau_- = \frac{1}{2}(\tau_x - i\tau_y),$$

where q is dimensionsless momentum $q = E_{(\nu,0)}b/\hbar c$. If the two-body spin-(isospin) operators

$$\mathcal{S}_\pm(i,j) = \frac{1}{2}(\sigma_i \pm \sigma_j),\quad T_\pm(i,j) = \frac{1}{2}(\tau_-(i) \pm \tau_-(j)), \qquad /12/$$

are defined, the space-like part (the time-like part is treated in analogy) is rewritten to give

$$\vec{J}_{(0)}(i) + \vec{J}_{(0)}(j) = \sum_{+,-}(j_1(qr_i)\hat{r}_i \pm j_1(qr_j)\hat{r}_j)\,(\mathcal{S}_+(i,j)\,T_\pm(i,j) + \qquad /13/$$

$$+ \mathcal{S}_-(i,j)\,T_-(i,j))$$

with $\hat{r}_i \equiv \frac{\vec{r}_i}{|\vec{r}_i|} = \sqrt{\frac{4\pi}{3}}\,Y^1(\hat{r}_i)$.

Starting with

$$i^L j_L(qr_i)\,Y^L_M(\hat{r}_i) = \frac{1}{4\pi}\int e^{i\vec{q}\cdot\vec{r}_i}\,Y^L_M(\hat{q})\,d\hat{q} \qquad /14/$$

and using the expansion of the plane wave into spherical harmonics the relation is obtained

$$j_L(kr_i)Y^L_M(\hat{r}_i) \pm j_L(kr_j)Y^L_M(\hat{r}_j) = \sum_{K,\ell}\sqrt{4\pi}\,(-)^{\frac{K+\ell+L}{2}}\,[1 \pm (-)^{\ell}]\,\times \qquad /15/$$

$$\times\, j_K(\bar{q}R)\,j_\ell(\bar{q}r)\,\hat{k}\,\hat{\ell}\,\hat{L}\begin{pmatrix}K & \ell & L\\ 0 & 0 & 0\end{pmatrix}\sum_{Q,m}(KQ\,\ell m|LM)\,Y^K_Q(\hat{R})\,Y^\ell_m(\hat{r}),\quad \bar{q} = q/\sqrt{2}.$$

Now, the n.m.e. of the one-body current becomes

$$\langle f|\vec{J}^{\,w}_{(0)}|i\rangle \sim \ell.c. \int_0^\infty R_{n\ell}(r)\,j_\ell(\bar{q}r)\,(1 - f(r))^2\,R_{n'\ell'}(r)\,r^2 dr. \qquad /16/$$

It is seen that in lowest order of eq./15/ and for low momenta ($j_0(\bar{q}r) \approx 1$) the change in the radial integrals is of the same magnitude as their deviation from the Kronecker symbol due to the correlation function $f(r)$, i.e., negligible. However, if operators of the type under consideration are involved in high-momentum processes, the effect of the repulsive two-body correlations can throughout become remarkable.

4. RESULTS AND CONCLUSIONS

In our previous calculation /3/ we have demonstrated that the effect of spreading of the 2p-2h admixtures over all possible $2\hbar\omega$-excitations for the 0^+_1-state and over the two-selected $3\hbar\omega$-excitations for the 0^-_1-state results in a destructive interference of the small 2p-2h contributions to the n.m.e. of the exchange current operator. This leads to the simple relation

$$_{1\hbar\omega + 3\hbar\omega}\langle 0^-_1;1|J^4_{(\pi')}|0^+;0\rangle_{0\hbar\omega+2\hbar\omega} \approx \alpha_0\beta_0\langle 2s1p^{-1}|J^4_{(\pi')}|0p-0h\rangle, \qquad /17/$$

α_0 and β_0 being the respective weights of the leading components of the initial and final nuclear states. According to this relation it is possible to illustrate clearly the sensitivity of the long-range and short-range pieces of $J^4_{(\pi')}$ to the short-range and medium polarization effects by studying only the right-hand side.

In the approximatly soft-current β-decay process the difference between the hard-pion and soft-pion results on the partial rate (the value of $\langle 0^-_1;1|J^{4(\beta^*)}_{(\rho A_1)}|0^+_1;0\rangle \equiv \langle J^{4(\beta^*)}_{(\rho A_1)}\rangle$) will measure the inconsistency in the determination of the transition operator and the nuclear wave functions. That's why once short-range correlations are included in the nuclear structure picture, the heavy-meson contribution reduces by almost of one order of magnitude. As a result, the two-body nuclear weak-axial charge density is strongly dominated (by 95%) by the one-pion exchange which is in accordance with chiral invariance statements. In the μ^--capture the short-range piece of the exchange operator behaves in a similar way (table 1).

The long-range part of the MEC operator in the hard pion model (eqs./1/, /3/) differs from the respective one-pion exchange operator $J^4_{(\pi)}$ of the soft pion model by the presence of the factors $\phi(y^2) \neq 1$. The latter indicate that a piece of $J^4_{(\pi)}$ originates from the $\rho\pi$- and $A_1\rho\pi$-weak decay currents. Because of that $J^4_{(\pi)}$ appears to be slightly more enhanced by the medium polarization effects ($\phi(m^{*2}_\pi) \approx 1.08$), and is much more damped by the hadronic form factors ($\phi(\Lambda^2_\pi) \approx 5$), on the other hand. The net effect including also the Jastrow correlations amounts in a reducement of the n.m.e. $\langle J^4_{(\pi)}\rangle$ by 40% compared to 20% for the exchange operator of the soft pion model, for which the value $\langle J^4_{(\pi)}\rangle = -0{,}1270$ is obtained.

Now, using the MEC operator from eq./3/ we shall carry out the shell-model calculation to get the ratio $\Lambda^\mu/\Lambda^\beta$ from which the ratio of the induced pseudoscalar to the axial-vector coupling constants g_p/g_A can than be extracted. In treating the one-body current we find that the simple relation /17/ remains valid only for the space-like part because it has the same (Gamow-Teller) form (eq./8/) as the MEC operator. The one-body weak axial

Table 1

Contributions to $<0_1^-, T=1|J^{4(\beta^*)}_{(\pi')}|0_1^+, T=0>$ with and without hadronic form factors (F), the correlation function f(r) and the medium polarization corrections (Π). $\Lambda^\beta_{(0)}(\Lambda^\beta_{corr})$ denotes the partial β-decay rate calculated without (with) correlation effects

Operator term	without f	/ with: f	f + F + Π
$<J^{4(\beta^*)}_{(\pi)}>$	-0.1511	-0.1129	-0.1055
$<J^{4(\beta^*)}_{(\rho A_1)}>$	0.0367	0.0085	0.0031
$<J^{4(\beta^*)}_{(\pi')}>$	-0.1144	-0.1044	-0.1024

Operator	$\bar{J}^4_{(\pi)}$	$J^4_{(\pi')}$
$\Lambda^\beta_{corr}/\Lambda^\beta_{(0)}$	0.75	0.95

charge density operator $(-j_0(kr_1)\vec{\sigma}_1\cdot\vec{\nabla}_1/M)$ to be renormalized by MEC effects is rather sensitive to the art of spreading of the 2p-2h-admixtures to the initial and final nuclear states. Because of that in calculating the one-body contribution to the n.m.e., the complete 70-component wave function of the 0_1^--state from ref. /11/ will be used. The results of the calculated μ^--capture and β-decay partial rates in this extended shell-model configuration space are given in table 2 for three different values of the oscillator strength parameter b. One can see from the table that the renormalization $\delta = <J^4_{(\pi)}>/<J^4_{(0)}>$ does not strongly depend on b. However, the absolute value of $\Lambda\beta$ changes by about 50% if b is varied by 10%. This is so because of the proportionality of the Gamow-Teller operators to b for small momenta $(j_1(x) \approx x/3)$. As a result, a value b = 1.6 fm, which differs from that extracted from electron scattering data, is likely to reproduce the data of Minamisono et al. /8/, whereas the theoretical prediction of Λ^β for b = 1.77 fm clearly favours the data of Gagliardi et al. /5/. The partial μ^--capture rate $(k_\mu^2 \approx 0.8\, m_\pi^2)$ is less sensitive to the b-value.

Table 2

The partial transition rates (in s^{-1}) calculated with the shell-model wave functions and the short-range plus medium polarization effects. The experimental value of the ratio of the transition rates $(\Lambda^\mu/\Lambda^\beta)_{exp} = (3.8 \pm 0.8)\cdot 10^3$ according to refs. /5,17/ or $(2.62 \pm 0.35)\cdot 10^3$ according to refs. /8,17/. The muon capture and beta-decay observables are denoted by μ and β respectively. For other notations see the text

n.m.e. b	1.6	1.7	1.77
$<J^{4(\beta^*)}_{(\pi')}>$	-0.0932	-0.0846	-0.0785
$<J^{4(\beta^*)}_{(0)}>$	-0.2336	-0.2198	-0.2111
δ_β	1.40	1.38	1.37
Λ^β	0.61	0.52	0.39
$<J^{4(\mu)}_{(\pi')}>$	-0.0824	-0.0730	-0.0674
$<J^{4(\mu)}_{(0)}>$	-0.1906	-0.1754	-0.1659
δ_μ	1.43	1.42	1.41
Λ^μ $(g_p/g_A = 13)$	1569	1420	1330
$\Lambda^\mu/\Lambda^\beta$	2572	2730	3410

So, the following conclusions can be drawn:

i) Account of repulsive two-body correlations of Jastrow type restore the dominance of the one-pion exchange in the two-body weak axial charge density operator of the hard-pion model which is in accordance with chiral-invariance statements.

Table 1

Contributions to $<0_1^-, T=1|J^{4(\beta^*)}_{(\pi')}|0_1^+, T=0>$ with and without hadronic form factors (F), the correlation function f(r) and the medium polarization corrections (Π). $\Lambda^\beta_{(0)}(\Lambda^\beta_{corr})$ denotes the partial β-decay rate calculated without (with) correlation effects

Operator term	without f	/ with: f	f + F + Π
$<J^{4(\beta^*)}_{(\pi)}>$	-0.1511	-0.1129	-0.1055
$<J^{4(\beta^*)}_{(\rho A_1)}>$	0.0367	0.0085	0.0031
$<J^{4(\beta^*)}_{(\pi')}>$	-0.1144	-0.1044	-0.1024

Operator	$\tilde{J}^4_{(\pi)}$	$J^4_{(\pi')}$
$\Lambda^\beta_{corr}/\Lambda^\beta_{(0)}$	0.75	0.95

charge density operator $(-j_0(kr_1)\vec{\sigma}_1\cdot\vec{\nabla}_1/M)$ to be renormalized by MEC effects is rather sensitive to the art of spreading of the 2p-2h-admixtures to the initial and final nuclear states. Because of that in calculating the one-body contribution to the n.m.e., the complete 70-component wave function of the 0_1^--state from ref. [11] will be used. The results of the calculated μ^--capture and β-decay partial rates in this extended shell-model configuration space are given in table 2 for three different values of the oscillator strength parameter b. One can see from the table that the renormalization $\delta = <J^4_{(\pi)}>/<J^4_{(0)}>$ does not strongly depend on b. However, the absolute value of $\Lambda\beta$ changes by about 50% if b is varied by 10%. This is so because of the proportionality of the Gamow-Teller operators to b for small momenta $(j_1(x) \approx x/3)$. As a result, a value b = 1.6 fm, which differs from that extracted from electron scattering data, is likely to reproduce the data of Minamisono et al. [6], whereas the theoretical prediction of Λ^β for b = 1.77 fm clearly favours the data of Gagliardi et al. [5]. The partial μ^--capture rate $(k^2_\mu \approx 0.8\, m^2_\pi)$ is less sensitive to the b-value.

Table 2

The partial transition rates (in s^{-1}) calculated with the shell-model wave functions and the short-range plus medium polarization effects. The experimental value of the ratio of the transition rates $(\Lambda^\mu/\Lambda^\beta)_{exp} = (3.8 \pm 0.8)\cdot 10^3$ according to refs. [5,17] or $(2.62 \pm 0.35)\cdot 10^3$ according to refs. [6,17]. The muon capture and beta-decay observables are denoted by μ and β respectively. For other notations see the text

n.m.e. b	1.6	1.7	1.77
$<J^{4(\beta^*)}_{(\pi')}>$	-0.0932	-0.0846	-0.0785
$<J^{4(\beta^*)}_{(0)}>$	-0.2336	-0.2198	-0.2111
δ_β	1.40	1.38	1.37
Λ^β	0.61	0.52	0.39
$<J^{4(\mu)}_{(\pi')}>$	-0.0824	-0.0730	-0.0674
$<J^{4(\mu)}_{(0)}>$	-0.1906	-0.1754	-0.1659
δ_μ	1.43	1.42	1.41
Λ^μ $(g_p/g_A = 13)$	1569	1420	1330
$\Lambda^\mu/\Lambda^\beta$	2572	2730	3410

So, the following conclusions can be drawn:

i) Account of repulsive two-body correlations of Jastrow type restore the dominance of the one-pion exchange in the two-body weak axial charge density operator of the hard-pion model which is in accordance with chiral-invariance statements.

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

Кирхбах М., Камалов С., Егер Х.-У. E4-84-37

Влияние корреляций на коротких расстояниях и поляризации среды на двухчастичную плотность слабого аксиального заряда в изовекторных переходах $^{16}O(0_1^+) \leftrightarrow {}^{16}N(0_1^-)$

Отталкивающие корреляции типа корреляций Джастрова в волновых функциях переходов $0_1^+ \leftrightarrow 0_1^-$, $\Delta T = 1$ в системе с $A = 16$ существенно подавляют вклад ρ- и A_1-мезона во временную компоненту обменного слабого аксиально-векторного тока в модели жестких пионов. Вместе с эффектами поляризации среды и адронных формфакторов они уменьшают вклад однопионного обмена на 40%. Применение маленького осцилляторного параметра ($b = 1{,}6$ Фм) позволяет описать одновременно экспериментальные величины парциальных скоростей μ^--захвата и β-распада /последние соответствуют данным японской группы/ для отношения константы индуцированного псевдоскалярного взаимодействия к аксиально-векторной константе $g_p/g_A = 13$.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

Сообщение Объединенного института ядерных исследований. Дубна 1984

Kirchbach M., Kamalov S., Jäger H.-U. E4-84-37

Short-Range and Medium Polarization Effects in the Two-Body Weak Axial Charge Density for the Isovector Transitions $^{16}O(0_1^+) \leftrightarrow {}^{16}N(0_1^-)$

Repulsive two-body correlations of the Jastrow type in the shell-model wave functions of the $0_1^+ \leftrightarrow 0_1^-$, $\Delta T = 1$ transitions in $A = 16$ nuclei are shown to damp crucially the ρ- and A_1-meson exchange contribution to the time component of the two-body weak axial vector current in the hard-pion model. Combined with medium polarization effects and hadronic form factors they reduce the one-pion exchange contribution by about 40%. With a small oscillator strength parameter ($b = 1.6$ fm) it is possible to reproduce simultaneously the experimental μ^--capture rate and the recent data in the partial β-decay rate reported by the Japan group for a ratio of the induced pseudoscalar to the axial-vector coupling constants $g_p/g_A = 13$.

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

Communication of the Joint Institute for Nuclear Research. Dubna 1984