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TWO-BODY WEAK AXIAL CHARGE DENSITY AND NUCLEAR STRUCTURE CORRELATION EFFECTS IN THE ISOVECTOR TRANSITION $^{16}O\,(0_1^+) \leftrightarrow {}^{16}N(0_1^-)$

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1. Within recent years the experimental determination of nuclear observables has achieved a high level of accuracy. This enables one to test current algebra and PCAC-predictions on weak coupling constant relations by analysing low-energy nuclear data. The accuracy reached by experimentalists is so that, in order to compare with theory, the theory must include correlations due to virtual meson exchange. The meson exchange in a nucleus provides not only the nuclear binding. Since the virtual mesons are charged, their transfer in the internucleon space gives rise to currents (mesonic exchange currents) which can interact with any external fields. So, all nuclear observables such as charge densities, magnetic moments or cross sections can be affected by the mesonic exchange. It also contributes to both the vector and axial-vector nuclear weak currents.

In the present paper we investigate the nuclear muon capture and beta decay processes between $^{16}O(0^+_1; T=0)$ and $^{16}N(0^-_1; T=1)$. As it was pointed out by Shapiro and Blokhintsev [1], the ratio $\Lambda_\mu/\Lambda_\beta$ of the partial muon capture rate $\Lambda_\mu(0^+\to 0^-)$ and the beta decay rate $\Lambda_\beta(0^-\to 0^+)$ tells us how the induced pseudoscalar nucleon form factor g_p is related to the axial nucleon form factor g_A. So, this purely axial first forbidden weak transition is well suited for testing experimentally the current algebra and PCAC statements which lead to $g_p/g_A \sim 7\text{-}8$. However, if the nuclear weak current is treated as a sum of the individual nucleon contributions, i.e., as a one-body current (this is known as the impulse approximation (IA)) the prediction of Λ_μ and Λ_β at the same level of accuracy becomes possible only if the ratio g_p/g_A is chosen to be $g_p/g_A \sim 13\text{-}20$. This is in contradiction to the current algebra estimation and appeares as a consequence of neglecting the mesonic exchange current. Indeed, the $0^+ \leftrightarrow 0^-$ transition is sensitive to the time component (the charge density) of the mesonic exchange axial-vector current, which is of the same order of magnitude $O(1/M)$ as the one-body current [2]. The axial-charge density in the isovector $0^+ \leftrightarrow 0^-$ transition in the A=16 system has recently been studied by several authors [3-6]. Some reservations, nevertheless, yet persist about the importance of the mesonic exchange contribution because of the simplified nuclear structure models [3-5] and the approximated operator form [3-6] used. The exchange operator was obtained from the S-matrix element of a Feynman diagram, describing only the long range one-pion exchange between two nucleons. The pion production amplitude by the weak current on the nucleon was derived from Adler's low energy theorems on pions with zero-momentum transferred (soft pions) [7].

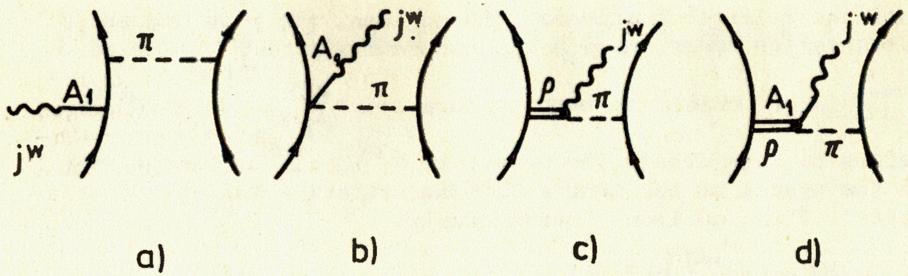
In the present investigation we improve the operator description by introducing short range (ρ and A_1 meson) exchanges as well as the nuclear structure model by spreading the 2p-2h-admixtures to the $0^+_1(0^-_1)$ states over $2\hbar\omega(3\hbar\omega)$ excitations in the $1s-(2p,1f)$space (in refs. [4,5] the 2p-2h admixtures were considered only to the groundstate of ^{16}O and the strength of the whole spectrum of the $2\hbar\omega$-excitations was assumed to be concentrated on only two selected components). We concentrate on the interplay between the mesonic exchange corrections and the nuclear-structure correlation effects.

2. The effects due to the presence of 2p-2h-admixtures to the nuclear states are calculated by using shell-model wave functions with configuration mixing. The latter were evaluated in ref. [8] by diagonalization of a nuclear residual interaction of Tabakin's type. The wave function of the ground state of ^{16}O contains all non-spurious $2\hbar\omega$-excitations. In the negative-parity state only the two strongest $3\hbar\omega$-excitations (~1% each) are preserved.

We also include in the nuclear structure picture nucleon-nucleon short range correlations by means of the correlation function $f(r)$ of Miller and Spencer (see ref. [9]).

$$f(r = r_i - r_j) = 1 - \exp(-\alpha r^2)(1 - \beta r^2),$$
$$\alpha = 1{,}1\ \mathrm{fm}^{-2}, \quad \beta = 0.68\ \mathrm{fm}^{-2}. \tag{1}$$

The weak axial-charge density operator is constructed as a two-body operator in the one-boson approximation by using the S-matrix method. As it was demonstrated by Weinberg [10] the low energy theorems can be derived from chiral invariant phenomenological Lagrangians by using only the lowest order zero-loop Feynman diagrams (tree-approximation). This idea enables one to treat consistently pions with a non-zero momentum transferred ("hard" pions) by extending the chiral SU(2)xSU(2) symmetry to a local gauge symmetry group and introducing vector and axial-vector mesons (ρ and A_1-mesons) as Yang-Mills fields. A minimal chiral invariant phenomenological Lagrangian for the $A_1\rho\pi$-system was proposed by Ogievetsky and Zupnik [11]. In order to calculate nuclear observables, this Lagrangian has been completed in ref. [12] to describe the $\Delta(1236)NA_1\rho\pi$-system. In the present paper the Feynman diagrams which contribute as O(1/M) (see the figure) are calculated with the phenomenological Lagrangian of the hard pion model from ref. [12]. So, we can treat correctly the currents generated by the ρ- and A_1-meson exchanges. After performing non-relativistic reduction, transformation to the coordi-



The Feynman diagrams which contribute as O(1/M) to the time part of the two-body exchange axial-vector isovector current in the hard pion model of ref. [12]. j^w stands for the weak leptonic axial-vector current.

nate space and multipole decomposition, the time component of the weak axial-vector current reads

$$J^{4(\mu,\beta)}_{(A_1\rho\pi)}(i,j) = \mp 4\pi\, T^{m=\pm 1}_{111}(r_i, r_j) \sum_{\substack{+,-;\\ C_\pm, A_\pm; \Gamma,\gamma;\\ K}} \mp i^{C_\pm + K} \sqrt{4\pi}\, Y^{*\Gamma}_{\gamma}(\hat{k}) \frac{\hat{C}_\pm \hat{K} \hat{A}_\pm}{\hat{\Gamma}\sqrt{3}} (-1)^K \times$$
$$\times (C_\pm 0 K 0 | \Gamma 0)(C_\pm 0 A_\pm 0 | 1 0) [\mathcal{F}^K(R) \otimes [\mathcal{L}^{A_\pm}(r) \otimes \mathcal{S}^1_\pm]^{C_\pm}]^\Gamma_\gamma \ . \tag{2a}$$

The following notations will be used:

$$\hat{L} \equiv \sqrt{2L+1}, \quad \vec{R} = \frac{1}{b\sqrt{2}}(\vec{r}_i + \vec{r}_j), \quad \vec{r} = \frac{1}{b\sqrt{2}}(\vec{r}_i - \vec{r}_j), \quad \hat{r} = \frac{\vec{r}}{|\vec{r}|},$$

$$\hat{k} = \frac{\vec{k}}{|\vec{k}|}, \quad C_+ = 2p, \quad A_+ = 2q+1, \quad C_- = 2p+1, \quad A_- = 2q,$$
$$p, q = 0, \dots, \infty,$$

$$T^{m}_{111}(r_i, r_j) = \sum_{m_1, m_2} (1m_1 1m_2 | 1m)\, r^1_{m_1}(i)\, r^1_{m_2}(j),$$

$$r^1_{\pm 1} = \mp \frac{1}{\sqrt{2}}(r_x \pm i r_y), \quad r^1_0 = r_z \ .$$

$j_\ell(x)$ is the regular spherical Bessel function, $\vec{r}_i$, $\vec{\sigma}_i$, τ_i refer to the position, spin, and isospin components of an i-th nucleon, k is the four-momentum associated with the axial current, g_r stands for the πNN coupling constant with $g_r/4\pi = 14{,}6$, f_π denotes the pion decay constant (f_π = 92 MeV), and b is the oscillator length parameter. The nucleon, the pion and the ρ meson masses are M, m_π and m_ρ, respectively. $J^{4(\mu)}_{(A_1\rho\pi)}$ with $T^{m=-1}_{111}(r_i,r_j)$ refers to muon capture and $J^{4(\beta)}_{(A_1\rho\pi)}$ with $-T^{m=1}_{111}(r_i,r_j)$ refers to beta decay. The operators $\mathcal{F}^K(R)$, $\mathcal{L}^{A\pm}(r)$ and $\mathcal{S}^1_\pm$ act in the spaces of the centre-of-mass, relative and spin coordinates of two nucleons, respectively:

$$\mathcal{F}^K_M(R) = j_K(Q^{(\mu,\beta)}R)\, Y^K_M(\hat{R}),$$

$$\mathcal{L}^{A\pm}_{\tilde{M}}(r) = j_{C_\pm}(Q^{(\mu,\beta)}r)\, Y^{A\pm}_{\tilde{M}}(\hat{r})(C^{(\pi)}\Phi^{(\pi)}(r) + C^{(A_1\rho\pi)}\,\Phi^{(A_1\rho\pi)}(r)),$$

$$\mathcal{S}^1_\pm = \vec{\sigma}^1_i \pm \vec{\sigma}^1_j, \quad Q^{(\mu,\beta)} = \frac{k^{(\mu,\beta)}}{\sqrt{2}}, \tag{2b}$$

$$k^{(\mu)} = \frac{E_\nu b}{\hbar c}, \quad k^{(\beta)} = \frac{E^\circ b}{\hbar c},$$

$$E_\nu = 95.121 \text{ MeV}, \quad E^\circ = 11.05 \text{ MeV}.$$

It is important to remark that the value of the oscillator length parameter b, extracted from the inelastic electron scattering data, depends strongly on the nuclear configuration mixing used. Here we use an averaged oscillator length $b = 1.7$ fm. $C^{(\pi)}$ and $\Phi^{(\pi)}(r)$ determine the common contribution from the πNN-pair and contact currents (graphs a,b)

$$C^{(\pi)} = -\frac{g_r^2 m_\pi^2}{16\pi g_A M^2},$$

$$\Phi^{(\pi)}(r) = Y_1(Am_\pi r), \tag{2c}$$

$$A = b\sqrt{2}, \quad Y_1(x) = \frac{e^{-x}}{x}(1 + \frac{1}{x}).$$

$C^{(A_1\rho\pi)}$ and $\Phi^{(A_1\rho\pi)}(r)$ determine the contribution from both the $\rho\pi$-weak decay and the $A_1\rho\pi$-currents (graphs c,d)

$$C^{(A_1\rho\pi)} = -\frac{g_r^2 m_\rho^2}{32\pi g_A M^2};$$

$$\Phi^{(A_1\rho\pi)}(r) = \int_0^1 dt\, e^{-Br} j_0(Akrt)\{1 + (\frac{1}{m_\rho^2} + \frac{\kappa_V}{4M^2})a^2[1 - \frac{2}{Br}(1 + \frac{1}{Br})] + \frac{\kappa_V}{4M^2 m_\rho^2} a^4[1 - \frac{4}{Br}(1 + \frac{1}{Br})]\}, \tag{2d}$$

with

$$a = \sqrt{t(1-t)k^2 + t(m_\rho^2 - m_\pi^2) + m_\pi^2},$$

$$\kappa_V = 3.7, \quad B \equiv ab\sqrt{2}.$$

In the soft current limit $k \to 0$ the function $\Phi^{(A_1\rho\pi)}(r)$ goes to

$$\Phi^{(A_1\rho\pi)}(r, k \to 0) \to F^{(\pi)} Y_1(Am_\pi r) + F^{(\rho A_1)} Y_1(Am_\rho r),$$

$$F^{(\pi)} \approx 2\frac{m_\pi^2}{m_\rho^2}, \tag{2e}$$

$$F^{(\rho A_1)} \approx -4(1 + \frac{\kappa_V m_\rho^2}{4M^2}).$$

Now the term $(C^{(\pi)} + C^{(A_1\rho\pi)}F^{(\pi)})Y_1(Am_\pi r)$ in $\mathcal{L}^{A\pm}(r)$ from eq. (2b) generates in eq. (2a) the standard representation of the long-range soft pion exchange operator $J^4_{(\pi)}$. The term $C^{(A_1\rho\pi)}F^{(\rho A_1)}Y_1(Am_\rho r)$ opposite in sign generates the short-range contribution from the ρ- and A_1-exchange, $\tilde{J}^4_{(\rho A_1)}$. Explicitly these contributions are identical to the well-known expressions [13]

$$J^4_{(\pi)} = -\frac{g_r^2 m_\pi^2}{8\pi g_A M^2} i(r_i \times r_j)_\pm (\vec{\sigma}_i + \vec{\sigma}_j)\cdot\hat{r}\cdot Y_1(Am_\pi r),$$

$$\tilde{J}^4_{(\rho A_1)} = \frac{g_r^2 m_\rho^2}{8\pi g_A M^2}(1 + \frac{\kappa_V m_\rho^2}{4M^2}) i(r_i \times r_j)_\pm (\vec{\sigma}_i + \vec{\sigma}_j)\cdot\hat{r}\cdot Y_1(Am_\rho r),$$

with $[\ ]_\pm = \frac{1}{2}([\ ]_x \pm i[\ ]_y)$.

So, we have

$$J^4_{(A_1\rho\pi)}(k \to 0) - J^4_{(\pi)} \to \tilde{J}^4_{(\rho A_1)}. \tag{2f}$$

3. The partial transition rates are proportional to the square of the matrix element of the weak axial current $J^{(\mu,\beta)}$ between the initial and final nuclear states

$$\Lambda_{(\mu,\beta)}(0^+_1 \rightleftarrows 0^-_1) \sim |<0^-_1, T=1|\, J^{(\mu,\beta^*)}|0^+_1, T=0>|^2 ,$$

$$J^{(\mu,\beta)} = \sum_{i=1}^{A} J^{(\mu,\beta)}_{IA}(i) + \sum_{i\neq j} J^{4(\mu,\beta)}_{(A_1\rho\pi)}(i,j)\ . \qquad (3)$$

Detailed expressions for $\Lambda_{(\mu,\beta)}$ and the one-body current are given in ref. [13]. In the second-quantization representation the matrix element of a spherical tensor component $J^{(\lambda,\bar\mu)}(t,r)$ of the current operator (λ, $\bar\mu$ and t, r are the respective multipolarities in the spin-coordinate and the isospin spaces) reads

$$<f|\, J^{(\lambda,\bar\mu;\, tr)}\, |i> = \sum_{a,b}{}' \frac{1}{\hat\lambda \hat t}(a||\, J^{(\lambda,t)}_{IA}||b)\,\rho^{i\to f}_{[ab]\lambda,\bar\mu;\, t,r} +$$

$$+ \sum_{\substack{a<b\\ c<d\\ J,T,J',T'}} \frac{1}{\hat\lambda \hat t}(ab;\, JT||\, J^{(\lambda,t)}_{(A_1\rho\pi)}||cd;\, JT)\,\rho^{i\to f}_{[ab]JT;\,[cd]J'T';\,\lambda,\bar\mu;\, t,r}\ . \qquad (4)$$

We apply Greek letters in order to denote the full set of quantum numbers defining a single-particle state. The corresponding Roman alphabet denotes the same set except for the magnetic quantum numbers

$$|\alpha> \equiv |n_a \ell_a j_a m_a;\ \tfrac{1}{2}\ r_a \ldots> \equiv |a,\ m_a, r_a>$$

The one- and two-body transition density matrices are defined as follows:

$$\rho^{i\to f}_{[ab]\lambda,\bar\mu;\, t,r} = <f|\ [a^+_a \otimes a_b]^{\lambda,\bar\mu}_{t,r}\ |i>,$$

$$\rho^{i\to f}_{[ab]JT;\,[cd]J'T';\lambda,\bar\mu;\, t,r} = (1+\delta_{ab})^{1/2}(1+\delta_{cd})^{1/2} <f|[[a^+_a a^+_b]^{JT} \otimes [a_c a_d]^{J'T'}]^{\lambda,\bar\mu}_{t,r}|i>. \qquad (5)$$

We use brackets to denote the tensor product of two tensor operators and the symbol $\otimes$ for the coupling

$$[B^+_{JT} \otimes B_{J'T'}]^{\lambda,\bar\mu}_{t,r} =$$

$$= \sum_{[M],[T_3]} (-)^{J'-M'+T'-T'_3}(JMJ'-M'|\lambda\bar\mu)(TT_3T'-T'_3|tr)\, B^+_{JM;TT_3} B_{J'M';\,T'T'_3}\ .$$

As far as many-body wave functions are concerned, we make use from the Rossendorf shell-model code [14] which was extended in order to generate two-body transition density matrices as required by eq. (5). The reduced two-body matrix elements from eq. (4) are calculated by first transforming to the L-S representation in order to separate the spin and spatial coordinates. Then the spatial matrix elements are transformed to relative and centre-of-mass coordinates using the Brody-Moshinsky coefficients (expressions are given in ref. [15]). Once the standard Racah algebra technique is applied and the integration over the angles is carried out, the reduced (coupled) two-body matrix elements of the exchange operator are given by a linear combination (Σ_q) of geometrical factors $\{\ldots\}$ and radial integrals of the type

$$(ab;\, JT||\, J^{(\lambda,t)}_{(A_1\rho\pi)}\, ||cd;\, J'T') \sim$$

$$\sim \Sigma_q \{\ldots\} \int_0^\infty R_{n\ell}(r)\,\Phi(r)\, j_{C_\pm}(Q^{(\mu,\beta)} r)\, R_{n'\ell'}(r)\, r^2 dr\ .$$

The dependence of the spatial wave function on the relative coordinate $R_{n\ell}(r)$ corresponds in our consideration to an oscillator state (the number of quanta n and angular momentum ℓ). Now the short-range correlations are introduced by the replacement $R_{n\ell}(r) \to f(r) R_{n\ell}(r)$.

4. To get an idea about the relevance of the mesonic exchange contributions, we consider first a rather simplified situation by assuming that the ground state of ^{16}O is a closed core ($|0p - 0h>$) and the 0^-_1 state is built up only of one particle-hole configuration $|(2s_{1/2})^1(1p_{1/2})^{-1} J=M=0;\ T=1, T=-1>$ (Table 1). In this case the transition density matrices look very simple

$$\rho^{0^+_1\to 0^-_1}_{[ab]0,0;\,1,-1} = \delta_{a,2s_{1/2}}\,\delta_{b,1p_{1/2}}\ ,$$

$$\rho^{0^+_1\to 0^-_1}_{[ab]JT;\,[cd]J'T';\,0,0;\,1,-1} = \qquad (6)$$

$$= n^{-1}_{1p_{1/2}d}\,\delta_{a,2s_{1/2}}\,\delta_{b,d}\,\delta_{c,1p_{1/2}}\,\frac{\hat J(-)^T}{2}\,\delta_{JJ'}\ ,$$

$$n_{cd} = (1-(-)^{J+T'}\delta_{c,d})^{-1/2}\ , \quad |b> = 1s_{1/2},\ 1p_{1/2},\ 1p_{3/2}\ .$$

Table 1

The partial transition rates for the muon capture and the beta-decay $\Lambda_{(\mu,\beta)}$(in s^{-1})calculated for $g_p/g_A = 7.5$ with the mesonic exchange corrections and the wave functions without configuration mixing. For a summary of the experimental situation see ref.[6]

Nuclear matrix element/partial transition rates	Muon capture		Beta decay	
$<0_1^-\|J_{IA}^{(\mu,\beta *)}\|0_1^+>$	-0.3246		-0.1098	
$\Lambda_{(\mu,\beta)}^{IA}$	2.7×10^3		0.40	
$\Lambda_{(\mu,\beta)}^{exp}$	$(1.57\pm0.1)\times10^3$		0.41 ± 0.06	
	without f	with f	without f	with f
$<0_1^-\|J_{(A_1\rho\pi)}^{4(\mu,\beta *)}\|0_1^+>$	-0.0986	-0.0896	-0.1144	-0.1044
$\Lambda_{(\mu,\beta)}$	3.74×10^3	3.63×10^3	1.02	0.96

Operator:	IA;	$J_{(A_1\rho\pi)}^4$ without f	with f	experiment
$\Lambda_\mu/\Lambda_\beta$	6.65×10^3	3.66×10^3	3.78×10^3	$(3.8\pm0.8)\times10^3$

A glance at Table 1 indicates that due to the mesonic exchange corrections the ratio $\Lambda_\mu/\Lambda_\beta$ is reduced strongly with respect to the impulse approximation. This is in accordance with the experimental data. The importance of accounting for short range correlations in the different pieces of the operator $J^4_{(A_1\rho\pi)}$ from eq.(2f) is illustrated on Table 1a for the beta decay process. One can see from the Table that the use of the hard pion exchange operator is almost equivalent to the use of the long range part $J^4_{(\pi)}$ with short range correlations considered. In the approximately soft-current beta decay process, the difference between the hard pion- and soft pion results (the value of $<\tilde{J}^4_{(\rho A_1)}>$) measures the inconsistancy in the determination of the transition operator and the nuclear wave functions. That's why, once the short range correlations are included in the operator (it contains ρ- and A_1-exchanges), as well as in the nuclear structure model (by means of the correlation function $f(r)$), the contribution $<\tilde{J}^4_{(\rho A_1)}>$ becomes negligible. The two-body axial charge density operator is strongly dominated by the soft pion exchange even if ρ- and A_1-exchanges are considered. So, the correct behaviour of the matrix elements of the hard pion operator constituents in the soft pion limit is ensured, once short range correlations are introduced in the nuclear structure description. The predicted values for the partial transition rates are, however, still too large. This is a consequence of the rather restricted configuration space. The results of the calculation with the many-body shell model wave functions with configuration mixing with and without short-range correlations are given in Table 2. Our calculations demonstrate the great importance of the 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 negative parity 0_1^- state. In this case - 900 non-vanishing reduced matrix elements

Table 1a

Contributions to $<0_1^-, T=1\|J^{4(\beta *)}_{(A_1\rho\pi)}\|0_1^+, T=0>$ with and without the correlation function $f(r)$ in eq. (1). $\Lambda_\beta^{(0)}$ ($\Lambda_\beta^{(f)}$) denote the partial beta-decay rate calculated without (with) $f(r)$. The hard (soft) pion model predictions are labelled by $A_1\rho\pi(\pi)$.

Operator term	without f	with f
$<J^{4(\beta *)}_{(\pi)}>$	-0.1511	-0.1129
$<\tilde{J}^{4(\beta *)}_{(\rho A_1)}>$	0.0367	0.0085
$<J^{4(\beta *)}_{(A_1\rho\pi)}>$	-0.1144	-0.1044
$\Lambda_\beta^{(A_1\rho\pi)}/\Lambda_\beta^{(\pi)}$	0.75	0.95
Operator type:	$J^4_{(\pi)}$	$J^4_{(A_1\rho\pi)}$
$\Lambda_\beta^{(f)}/\Lambda_\beta^{(0)}$	0.79	0.94

Table 2

The partial transition rates calculated with the shell-model wave function with configuration mixing and including mesonic exchange corrections (in s^{-1})

g_p/g_A	7.5		10.5	
Λ_μ^{IA}	1.8×10^8		1.3×10^8	
Λ_β^{IA}	0.18		0.18	
	without f	with f	without f	with f
Λ_μ	2.5×10^8	2.43×10^8	1.9×10^8	1.84×10^8
Λ_β	0.53	0.50	0.53	0.50
$\Lambda_\mu/\Lambda_\beta$	4.7×10^8	4.86×10^8	3.6×10^8	3.68×10^8

$(a,b;JT\|J^{(0,1)}_{(A_1\rho\pi)}\|cd;JT')$ (see eq. (4)) contribute to the nuclear matrix element in eq. (3) (compared to 8 matrix elements in the calculation without the configuration mixing). As a consequence of the destructive interference of the small contributions it becomes evident that the 2p-2h-admixtures affect the mesonic exchange corrections only through the change in the weights of the leading components of the nuclear wave functions. So, the nuclear matrix elements of the exchange operators

$$<0_1^-, T=1\,|\,J^{4(\mu)}_{(A_1\rho\pi)}\,|\,0_1^+, T=0> \quad \text{and} \quad <0_1^-, T=1\,|\,J^{4(\beta*)}_{(A_1\rho\pi)}\,|\,0_1^+, T=0>$$

in Table 2 are smaller than the corresponding matrix elements from Table 1 approximately by a factor of $R=a_0\beta_0$, where a_0 and β_0 are the respective weights of the leading components $|0p\text{-}0h>$ and $|(2s\,{}^1\!/_2)^1(1p\,{}^1\!/_2)^1>$ of the 0_1^+ and 0_1^- states ($a_0\approx0.89$, $\beta_0=0.95$). The nucleon-nucleon short-range correlations reduce additionally the nuclear matrix elements of the hard-pion exchange operator by about 10% (Tables 1 and 2). The absolute values of the partial transition rates Λ_μ and Λ_β change, however, by only 3% and 6%, respectively.

5. So, the conclusion can be made that the mesonic exchange corrections do not dramatically depend on the details of the realistic nuclear-structure picture and that they are very important for the interpretation of the experimental data on nuclear weak axial-charge density phenomena. A further glance on Table 2 indicates that for the value of $g_p/g_A\sim7$, as it is predicted by current algebra relations, the main bulk of the partial muon capture rate lies in the impulse approximation. The choice of $g_p/g_A\sim10$ enables one to predict Λ_μ at the same level of accuracy ($\sim$20%) as Λ_β. So, some indication becomes evident that the ratio g_p/g_A may slightly deviate from the current algebra prediction.

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Егер Х.-У., Кирхбах М., Труглик Э. E4-83-130

Мезонные обменные поправки к плотности аксиального заряда и их связь с ядерной структурой в изовекторном переходе $^{16}O(0_1^+) \leftrightarrow {}^{16}N(0_1^-)$

Показано, что улучшение модели механизма реакции $^{16}O(0_1^+) \underset{\text{бета-распад}}{\overset{\mu\text{ захват}}{\rightleftarrows}} {}^{16}N(0_1^-)$ путем учета обмена пионами и векторными ρ- и A_1-мезонами в модели жестких пионов и улучшение модели ядерной структуры путем учета корреляционных эффектов приводит к такому уровню теоретического описания процессов, который позволяет из ядерно-физического эксперимента извлечь информацию о фундаментальных константах слабого взаимодействия. В результате для отношения индуцированного псевдоскалярного формфактора g_p к аксиальному формфактору нуклона g_A получается величина $g_p/g_A \sim 10$, которая довольно близка к предсказанию алгебры токов ($g_p/g_A \sim 7$-8). Учет корреляций на коротких расстояниях в ядерных волновых функциях необходим для обеспечения правильного поведения результатов, полученных методом жестких пионов, в низкоэнергетическом пределе.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

Препринт Объединенного института ядерных исследований. Дубна 1983

Jäger H.-U., Kirchbach M., Truhlik E. E4-83-130

Two-Body Weak Axial Charge Density and Nuclear Structure Correlation Effects in the Isovector Transition $^{16}O(0_1^+) \leftrightarrow {}^{16}N(0_1^-)$

Hard pion model description of mesonic exchange corrections to the weak processes $^{16}O(0_1^+) \underset{\text{beta decay}}{\overset{\text{muon capture}}{\rightleftarrows}} {}^{16}N(0_1^-)$ as well as nuclear structure correlation effects lead to a level of accuracy of the theoretical data analysis which enables one to obtain information about weak coupling constants. So, the extracted ratio of the induced pseudoscalar coupling constant g_p to the axial nucleon form factor g_A ($g_p/g_A \sim 10$) lies close to the current algebra prediction ($g_p/g_A \sim 7$-8). Short range correlation effects, introduced in the nuclear wave functions, are necessary because they guarantee a correct behaviour of the hard pion results in the soft pion limit.

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

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