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MESON EXCHANGE CORRECTIONS TO NUCLEAR WEAK AXIAL CHARGE DENSITY IN HARD PION MODEL AND $0^+ \leftrightarrow 0^-$ TRANSITION IN A=16 NUCLEI

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1. INTRODUCTION

The two-body part of the time-like component of the weak axial-vector current attracts attention since Kubodera, Delorme and Rho [1] have pointed out that it is of the same magnitude $O(1/M)$ as the one-body part. Due to the approximate chiral invariance of the strongly interacting nucleon system the two-body part (the mesonic exchange correction (MEC)) could give rise to large effects even in the one-pion exchange (OPE) limit. This can be proved by calculating the nuclear axial charge density. For this aim the purely axial weak processes of the beta decay and the muon capture between $^{16}N(0^-_1, T=1)$ and $^{16}O(0^+_1, T=0)$ are well suited because they are rather sensitive to the time component of the current. The ratio $\Lambda_\mu/\Lambda_\beta$ of the partial muon capture rate $\Lambda_\mu(0^+ \to 0^-)$ to the partial beta decay rate $\Lambda_\beta(0^- \to 0^+)$ gives us information how the induced pseudoscalar coupling constant g_p is related to the axial nucleon form-factor g_A(ref. [2]). The importance of accounting for MEC in this first forbidden transition has been discussed in a series of papers [3-8]. The most extensive investigation was performed by Towner and Khanna [8]. They estimated for the soft pion MEC operator [1] the operator renormalization for a variety of residual interactions by second-order perturbation theory along the standard line of the linked valence cluster expansion of Brandow. For one residual interaction (OBEP) out of six they reproduced the experimental values for Λ_μ and Λ_β only with meson exchange corrections incorporated in the transition operator for $g_p/g_A = 10$. Here we concentrate our attention on the form of the two-body operator, which is of more complicated structure than it is predicted by soft-pion low energy theorems. In order to construct the two-body current operator we exploit the phenomenological Lagrangian (PL) version [9,10] of the hard pion model [11,12]. Further we investigate the interplay between the heavy meson exchanges and the nuclear structure correlation effects. We calculate directly the effects from the 2p-2h admixtures to the 0^+ and 0^- nuclear states using shell model wave functions with configuration mixing [13]. The latter are evaluated by diagonalization of a residual interaction due to Tabakin. The 0^+ wavefunction contains all possible non-spurious $2\hbar\omega$-excitations whereas the 0^- wavefunction includes only selected $3\hbar\omega$-excitations.

The two nucleon MEC operator is developed in sec.2. The partial transition rates are calculated in sec.3. Concluding remarks are given in sec.4.

2. THE TWO-BODY AXIAL MEC OPERATOR

Constructing the two-nucleon MEC operator in the one-boson (tree-) approximation we start with the hard pion model for the $A_1\rho\pi$-system as proposed by Ogievetsky and Zupnik [9] which was extended in ref. [10] to the $N\Delta(1236)A_1\rho\pi$-system in order to describe the nuclear muon capture and beta decay. The PL of the hard pion model is chosen to be invariant under the local $SU_2 \times SU_2$ transformation. This chiral gauge invariance is assumed to be broken only due to the non-zero masses m_ρ and m_{A_1} of the ρ - and A_1-mesons, respectively. The A_1-meson guarantees the consistency of the chiral approach with vector dominance [11,12]. The PL of the hard pion model provides a good description of hadron amplitudes up to ~ 1 GeV already in the tree approximation. In the Ogievetsky-Zupnik version of the hard pion model the effective PL for the $A_1\rho\pi$-system is completely determined by four phenomenological parameters: the masses of the respective mesons and their coupling constants (g_ρ, g_{A_1}). In this way an unambiguous counting of the pion and heavy meson exchange graphs is ensured. This is an advantage compared to the method of current algebra (CA) and partial conserved axial-vector current (PCAC) where only the pion coupling constants are fixed. The PL technique is convenient for practical calculations, because it enables one to apply the standard Feynman rules. We define the two-nucleon MEC operator as a set of all possible tree-graphs which are required for a given process. The resulting operator possesses the correct chiral $SU_2 \times SU_2$ transformation properties and reproduces all standard PCAC results in the soft pion limit. The corresponding diagrams which contribute significantly to the transition rates are displayed in fig.1. In the subsequent calculations the space component of the exchange current will be completely neglected because it is by one order of magnitude smaller ($O(1/M^2)$) than its time component ($O(1/M)$). The space component of the nuclear current is considered only in the one-body part (i.e., in the impulse approximation (IA)) in the standard way [14]. In the non-relativistic limit and after performing the transformation to the coordinate space (cf. the work of Chemtob and Rho [15]) the operator of the nuclear weak axial charge density reads

$$J^{4(\mu,\beta)}_{\rm exch} = \sum_{i\neq j} e^{-i\vec{k}\cdot\vec{r}_i}\, \vec{\sigma}_i\cdot\hat{r}_{ij}\, \{C_1 Y_1(A m_\pi r) + C_2 F_0(r)\}\, i\,(\vec{\tau}_i\times\vec{\tau}_j)_{\mp},$$

$$Y_1(x) = \frac{e^{-x}}{x}\left(1+\frac{1}{x}\right),$$

$$\vec{r} = \frac{1}{b\sqrt{2}}(\vec{r}_i - \vec{r}_j), \qquad \hat{r}_{ij} = \frac{\vec{r}}{|\vec{r}|},$$

$$C_1 = -\frac{g_r^2 m_\pi^2}{8\pi g_A M^2}\left[\frac{1}{2} + \frac{2 f_{\pi NN^*} f_\pi}{9M(M^*-M)}\right],$$

$$C_2 = -\frac{g_r^2 m_\rho^2}{32\pi g_A M^2}, \qquad A = b\sqrt{2}. \tag{2.1}$$

Here $\vec{r}_i$, $\vec{\sigma}_i$, $\vec{\tau}_i$ refer to the position, spin, and isospin components of the i-th nucleon, k is the four-momentum associated with the axial current, g_r ($f_{\pi NN^*}$) stand for the respective πNN(πNN^*) coupling constants with $g_r^2/4\pi = 14.6$ ($f^2_{\pi NN^*}/4\pi = 0.23$) and f_π denotes the pion decay constant ($f_\pi = 92$ MeV). The $\Delta(1236)$-isobar, the nucleon and the pion masses are M^*, M and m_π, respectively, and b is the oscillator length parameter. $J^{4(\mu)}_{\rm exch}$ with $(\vec{\tau}_i\times\vec{\tau}_j)_-$ refers to muon capture and $J^{4(\beta)}_{\rm exch}$ with $(\vec{\tau}_i\times\vec{\tau}_j)_+$ refers to beta decay. The function $F_0(r)$ has the explicit form

$$F_0(r) = \int_0^1 dt\, e^{-Br} j_0(Akrt)\, \Big\{1 + \Big(\frac{1}{m_\rho^2} + \frac{\kappa_\nu}{4M^2}\Big) a^2 \Big[1 - \frac{2}{Br}\Big(1+\frac{1}{Br}\Big)\Big] + \frac{\kappa_\nu}{4M^2 m_\rho^2} a^4 \Big[1 - \frac{4}{Br}\Big(1+\frac{1}{Br}\Big)\Big]\Big\} \tag{2.2}$$

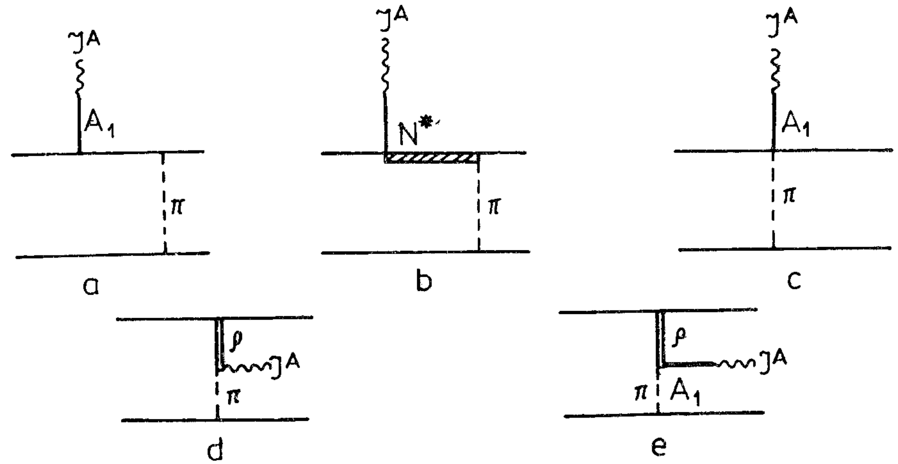

Fig.1. Feynman graph representation of the two-nucleon axial MEC operator in the tree-approximation (see ref. [10]): a) pair term; b) isobar excitation current; c) constant term; d) $\rho\pi$-weak decay current; e) $A_1\rho\pi$-current; J^A stands for the weak leptonic axial-vector current.

with

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

$$\kappa_v = \mu_p - \mu_n = 3.7, \qquad B = ab\sqrt{2},$$

and determine the radial dependence of the contributions from both the $\rho\pi$-weak decay current and the $A_1\rho\pi$-current (graphs d,e, in fig.1). The correct treatment of $F_0(r)$ is necessary in the muon capture because of the large four-momentum transferred ($k \approx 0.8 m_\pi^2$). The term with the denominator $1/M(M^*-M)$ represents the small contribution coming from the Δ-isobar and will be neglected. In the soft current limit ($k \to 0$) and for $m_\pi/m_\rho \ll 1$, $m_\pi/M \ll 1$ in $F_0(r)$ the exchange operator of the hard pion model (2.1) can be separated into two parts. The first one represents nothing else than the MEC operator developed from the low energy theorems /1,3,7/. From now on the symbol $J^{4(\mu,\beta)}_{OPE}$ will be used in order to denote the one pion exchange operator. In doing this we obtain

$$J^{4(\mu,\beta)}_{exch}\Big|_{k\to 0} = J^{4(\mu,\beta)}_{OPE} + \sum_{i\neq j} \ell^{i\vec{k}\cdot\vec{r}_i}\, \vec{\sigma}_i\cdot\hat{\vec{r}}_{ij}\, C_2'\, Y_1(Am_\rho r)\, i(\vec{\tau}_i \times \vec{\tau}_j)_\mp ,$$

$$J^{4(\mu,\beta)}_{OPE} = -\frac{g_r^2 m_\pi^2}{8\pi g_A M^2} \sum_{i\neq j} \ell^{i\vec{k}\cdot\vec{r}_i}\, \vec{\sigma}_i\cdot\hat{r}_{ij}\, Y_1(Am_\pi r)\, i(\vec{\tau}_i\times\vec{\tau}_j)_\mp , \qquad (2.3)$$

$$C_2' = \frac{g_r^2 m_\rho^2}{8\pi\, g_A M^2}\left(1 + \frac{\kappa_v m_\rho^2}{4M^2}\right).$$

It should be noted that the term C_2' is larger by a factor $(1+\kappa_v m_\rho^2/4M^2) \sim 1.6$ as compared to ref. /7/ where only the $\rho\pi$-exchange graph was considered. This is due to the contribution coming from the coupling of the ρ-meson to the anomalous magnetic moment of the nucleon.

3. RESULTS FOR THE MESONIC EXCHANGE CORRECTIONS TO THE TRANSITION RATES

We consider the muon capture reaction

$$\mu^- + {}^{16}O\,(0_1^+;\, T=0) \to {}^{16}N\,(0_1^-;\, T=1) + \nu_\mu$$

and the inverse reaction of beta decay

$${}^{16}N\,(0_1^-;\, T=1) \to {}^{16}O\,(0_1^+;\, T=0) + e^- + \bar{\nu}_e .$$

The partial transition rates are given by /7/

$$\Lambda_\mu(0^+\to 0^-) = 8.375(\sqrt{2}\, g_A(k^2))^2 \times 10^3 |\langle 0^-|\, J^{(\mu)}\, |0^+\rangle|^2, \qquad (3.1a)$$

$$g_A(k^2) = 1.24,$$

$$\Lambda_\beta(0^-\to 0^+) = 10.65\,(\sqrt{2}\, g_A(0))^2 |\langle 0^-|\, J^{(\beta)*}\, |0^+\rangle|^2, \qquad (3.1b)$$

$$g_A(0) = 1.26 .$$

The current operators $J^{(\mu)}$, $J^{(\beta)*}$ which include the mesonic exchange corrections have the form

$$J^{(\mu,\beta^*)} = J^{(\mu,\beta^*)}_{IA} + \frac{1}{\sqrt{2}\, g_A} J^{(\mu,\beta^*)}_{exch} . \qquad (3.2a)$$

The one-body operators $J^{(\mu,\beta^*)}_{IA}$ are kept in mind as

$$J^{(\mu)}_{IA}(i) = \frac{1}{M} j_0(k^{(\mu)} r_i)\vec{\sigma}_i\cdot\vec{\nabla}_i\,(\vec{\tau}_-)_i - g^{(\mu)} j_1(k^{(\mu)} r_i)\vec{\sigma}_i\cdot\hat{r}_i\,(\vec{\tau}_-)_i$$

$$J^{(\beta)*}_{IA}(i) = \frac{1}{M} j_0(k^{(\beta)} r_i)\vec{\sigma}_i\cdot\vec{\nabla}_i\,(\vec{\tau}_-)_i + g^{(\beta)} j_1(k^{(\beta)} r_i)\vec{\sigma}_i\cdot\hat{r}_i\,(\vec{\tau}_-)_i \qquad (3.2b)$$

with

$$g^{(\mu)} = \left(1 - \left(\frac{g_p}{g_A(k^2)} - 1\right)\frac{E_\nu}{2M}\right), \qquad g_p/g_A = 7.5,$$

$$g^{(\beta)} = 0.932\left(1 + \frac{3\alpha Z}{2RE^\circ}\right), \qquad R = 3.51\ \text{fm}, \quad \alpha = \frac{1}{137},$$

$$k^{(\mu)} = E_\nu b/\hbar c, \quad k^{(\beta)} = E^0 b/\hbar c, \quad E_\nu = 95.121\ \text{MeV}, \quad E_0 = 11.05\ \text{MeV}, \ b = 1.7\ \text{fm}.$$

In order to obtain information about the relevance of the heavy meson exchange contributions we consider first a rather simplified situation by assuming that the 0^+ state in a closed core and the 0^- state is built up by only one particle-hole configuration $|(2s_{1/2})^1(1p_{1/2})^{-1};\ J=M=0;\ T=1, T_3=-1\rangle$. In this case the nuclear matrix element reads

$$\langle 0^- | J^{(\mu,\beta^*)} | 0^+ \rangle =$$

$$= \frac{1}{\sqrt{3}} (2s_{1/2} || J_{1A}^{(\mu,\beta^*)} || 1p_{1/2}) + \frac{1}{\sqrt{2}\, g_A} \times \tag{3.3}$$

$$\times \frac{1}{2\sqrt{3}} \sum_{\substack{n\ell j = 1s_{1/2}, \\ 1p_{1/2}, 1p_{3/2} \\ J = 0,1,2; \\ T,T' = 0,1}} \sqrt{1-\delta(n\ell j)(1p_{1/2})}\, (-)^{J+T'} \sqrt{2J+1}\, (-)^{T} (2s_{1/2}; n\ell j; JT || J_{exch}^{4(\mu,\beta^*)} || 1p_{1/2}; n\ell j; JT').$$

We define the renormalization of the nuclear matrix element as

$$\delta_{(\mu,\beta)} = 1 + \frac{1}{\sqrt{2}\, g_A} \frac{\langle 0^- | J_{exch}^{4(\mu,\beta^*)} | 0^+ \rangle}{\langle 0^- | J_{1A}^{(\mu,\beta^*)} | 0^+ \rangle}, \tag{3.4}$$

where the quantity

$$\theta = \frac{\delta_\mu^2}{\delta_\beta^2} = \frac{\Lambda_\mu \Lambda_\beta^{1A}}{\Lambda_\beta \Lambda_\mu^{1A}} \tag{3.5}$$

determines the reduction of the ratio $\Lambda_\mu/\Lambda_\beta$ compared to the value $\Lambda_\mu^{1A}/\Lambda_\beta^{1A}$ calculated without exchange effects. The corresponding results are given in table 1. It can be seen from the table that the partial rates derived from the hard pion operator $J_{exch}^{4(\mu,\beta)}$ are smaller than those obtained by the standard version $J_{OPE}^{4(\mu,\beta)}$ (compare ref. /3/). This results because the heavy meson graphs lead to an additional term (C_2' from (2.3)) which is of sign opposite to $J_{OPE}^{4(\mu,\beta)}$. Due to the contributions from the $\rho\pi$ - and $A_1\rho\pi$-diagrams the quantity θ changes in the right direction. The experimental value of the ratio $\Lambda_\mu/\Lambda_\beta$ is well reproduced, however, the predicted partial transition rates are still too large. This is a consequence of the rather restricted configuration space.

Now we shall improve the description of the nuclear structure by using shell model wavefunctions with configuration mixing. The relevant nuclear matrix elements are calculated applying the hard pion operator (2.1) and nuclear wavefunctions from ref. /13/ in which 2p-2h-admixtures are explicitly present. The groundstate wavefunction of ^{16}O was obtained by diagonalisation of a non-local separable potential of Tabakin in the whole space of all non-spurious $(0\hbar\omega, 2\hbar\omega)$ configurations. The wavefunctions of the negative parity states contain some uncertainties. A straightforward diagonalization of the Hamiltonian

Table 1

The renormalization of the single-particle matrix element $(2s_{1/2} - 1p_{1/2})$ due to mesonic exchange corrections (eq. (3.3)) both for the muon capture and the beta decay. The reduced matrix elements (r.m.e.) of the two-body exchange operator are labelled by $((n\ell j)_1 (n\ell j)_2;\ JT || J^{4(\mu,\beta^*)} || (n\ell j)_3 (n\ell j)_4;\ JT')$. For other symbols see the text. Experimental rates are $\Lambda_\mu = (1.57 \pm 0.10) \times 10^3 s^{-1}$ (ref. /16/) and $\Lambda_\beta = 0.41 \pm 0.06\ s^{-1}$ (ref. /17/)

quantum numbers of the r.m.e.		muon capture r.m.e.		beta decay r.m.e.	
$n\ell j$	JTT′	$J_{OPE}^{4(\mu)}$	$J_{exch}^{4(\mu)}$	$J_{OPE}^{4(\beta^*)}$	$J_{exch}^{4(\beta^*)}$
$2s_{1/2}\ 1s_{1/2}\ 1p_{1/2}\ 1s_{1/2}$	101	-.0593	-.0448	-.0592	-.0445
	110	.0586	.0489	.0698	.0606
$2s_{1/2}\ 1p_{1/2}\ 1p_{1/2}\ 1p_{1/2}$	001	-.0007	.0018	.0086	.0130
	110	.0432	.0304	.0557	.0437
$2s_{1/2}\ 1p_{3/2}\ 1p_{1/2}\ 1p_{3/2}$	110	.0219	.0118	.0240	.0161
	101	-.0003	.0008	.0037	.0056
	210	.0500	.0392	.0677	.0553
	201	-.0004	-.0002	.0048	.0072
$\langle 0^- \| J^{4(\mu,\beta^*)} \| 0^+ \rangle$		-.1335	-.0986	-.1511	-.1144
$\langle 0^- \| J_{1A}^{(\mu,\beta^*)} \| 0^+ \rangle$		-.3246		-.1098	
δ		1.23	1.17	1.77	1.58
$\Lambda_{(\mu,\beta)} (s^{-1})$		4.14×10^3	3.74×10^3	1.28	1.02
Type of the operator:		$J_{OPE}^{4(\mu,\beta^*)}$	$J_{exch}^{4(\mu,\beta^*)}$	experiment	
θ		0.48	0.55	(0.57)	

matrix for the 0^- state in the complete ($1\hbar\omega, 3\hbar\omega$) configuration space would require about 1800 several components of 2p–2h and 3p–3h type. In ref./13/ several wavefunctions of negative parity states were built up upon much more restricted model space. They describe correctly, e.g., the integral transition rates of the (μ, ν_μ) and (π^-, γ) reactions. As for the partial transition rates to the low lying negative parity states they are still too collective. The predicted values for $\Lambda_\mu(0^+ \to f)$ in the IA are by a factor 1.5-2 larger than the measured one (for details see ref./18/). In spite of this we use the nuclear structure model proposed in ref./19/ in order to investigate the possible consequences of $3\hbar\omega$- admixtures to the 0^-_1 (120 keV) state. In the wavefunction we preserve the admixtures (about 1% each) of two $3\hbar\omega$-configurations which come the strongest (normalization of the function is indeed that of ref./13/).

$$|^{16}N; 0^-_1, T=1\rangle = 0.9505\,|(2s_{1/2})^1(1p_{1/2})^{-1}J=0, T=1\rangle + $$
$$+ 0.0158\,|(1d_{3/2})^1(1p_{3/2})^{-1}J=0,\ T=1\rangle \qquad (3.6)$$
$$-0.1108\,|[(1d_{5/2})^1(1f_{5/2})^1\ J_p=0, T_p=0;\ (1p_{1/2})^{-2}\ J_h=0, T_h=1]\ J=0, T=1\rangle$$
$$-0.1036\,|[(2s_{1/2})^1\ (1f_{7/2})^1\ J_p=3,\ T_p=1;\ (1p_{3/2})^{-2}\ J_h=3,\ T_h=0]\ J=0, T=1\rangle.$$

In the model subspace spanned over the $1s_{1/2}$ up to $2p_{1/2}$ oscillator shells the J^4_{exch} operator is determined by 975 various non-vanishing reduced matrix elements in the jj -coupling scheme. Due to the spreading of the 2p–2h admixtures over all possible $2\hbar\omega$-excitations for the 0^+ state and over the two selected $3\hbar\omega$-excitations (eq. (3.6)) for the 0^- state the most important matrix elements (900) enter the calculation*. The result of this extended calculation of the transition rates $\Lambda_\mu(0^+ \to 0^-)$ and $\Lambda_\beta(0^- \to 0^+)$ is given in table 2. For the nuclear matrix elements of the exchange operator (2.1) we obtain

$$\langle 0^- |\ J^{4(\mu)}_{exch}\ | 0^+\rangle = -0.0812$$
$$\langle 0^- |\ J^{4(\beta)*}_{exch}\ | 0^+\rangle = -0.0945.$$

Compared to the calculation without configuration mixing (table 1) the matrix element of the two-body operator between the initial and final nuclear states reduces approximately by a factor of $R = \alpha_0 \beta_0$, where α_0 and β_0 are the respective weights of the leading components ($|0p - 0h\rangle$, $|(2s_{1/2}(1p_{1/2})^{-1}\ J=0, T=1\rangle$) of the 0^+ and 0^- states ($\alpha_0 = 0.89$, $\beta_0 = 0.95$). In other words it

*In the calculation without configuration mixing only 8 matrix elements contribute.

Table 2

The partial transition rates (in s^{-1}) calculated with the shell model wavefunctions with configuration mixing for two different values of the ratio g_p/g_A. Λ_μ and Λ_β stand for the partial rates evaluated using the exchange operator of the hard pion model (see eq. (2.1)) for the muon capture and the beta decay, respectively. The experimental value of the ratio of the transition rates is $(\Lambda_\mu/\Lambda_\beta)_{exp} = (3.8 \pm 0.8)\times 10^3$

g_p/g_A	7.5	10.5
$\Lambda^{IA}_\mu(0^+ \to 0^-)$	1.8×10^3	1.3×10^3
$\Lambda^{IA}_\beta(0^- \to 0^+)$	0.18	0.18
$\Lambda_\mu(0^+ \to 0^-)$	2.5×10^3	1.9×10^3
$\Lambda_\beta(0^- \to 0^+)$	0.53	0.53
δ_μ	1.17	1.20
$\dfrac{\Lambda_\mu(0^+ \to 0^-)}{\Lambda_\beta(0^- \to 0^+)}$	4.7×10^3	3.7×10^3

means that the main effect from the higher configurations is due to the change in the weights of the closed shell configuration of ^{16}O and of the dominating 1p–1h configuration in the 0^- state in ^{16}N. A similar result was found in ref./8/. However, this simple picture does not take place for the one-body part of the current (14 single particle matrix elements in the model subspace of consideration). As discussed above the predicted values for the partial transition rates in the IA appear to be overestimated. This is one of the reasons why the partial transition rates Λ_μ and Λ_β calculated in the present investigation are somewhat larger than the reported data. Other changes in transition rates are to be expected from finite size effects/6/ as well as from dependence of the πNN vertex on pion momentum in the two-body operator.

4. SUMMARY AND CONCLUDING REMARKS

In the present paper we studied in detail the mesonic exchange corrections in the $0^+ \leftrightarrow 0^-$, $\Delta T = 1$ weak axial transition in A = 16 nuclei by using shell model wavefunctions with configuration mixing. The results are as follows:

i) We have found that the MEC effect depends less on details of the nuclear wavefunction than the impulse approximation results.

ii) The main contribution to MEC effect is due to OPE part of the two-body operator (table 1). However, the heavy meson exchanges (graph 1d,1e) cannot be neglected in detailed calculations. Their influence is of the same order (~20%) as that expected from other corrections.

iii) The MEC effect is important for describing the nuclear weak axial charge density (table 2).

These results support strongly the concept presented and advocated in recent paper [19] by Guichon and Samour. As discussed by them the reason why MEC are expected to play an essential role in the $0^+ \to 0^-$ transition is of a principal nature and reflects the approximate global chiral symmetry properties of the strongly interacting nucleon system. The chiral invariance of a system consisting of non-zero mass particles can be realized only via the mechanism of spontaneous symmetry breaking, which leads to the appearance of massless Goldstone particles which could be, with a good accuracy, identified with pions. If the philosophy of chiral invariance of strong interaction is true and the mechanism of spontaneous symmetry breaking takes place, one should not expect, that the generator of axial chiral rotation (the axial charge) can be obtained as a simple sum of the purely nucleonic contributions with the pionic mode completely neglected, as impulse approximation suggests.

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APPENDIX

The Reduced Matrix Elements of the Weak Axial Charge Density MEC Operator

The separation of the relative coordinates from the centre-of-mass coordinates in (2.1) leads to the following expression for the time-like component of the weak axial-vector MEC operator:

$$J^{4(\mu,\beta^*)}_{\mathrm{exch}} = \frac{1}{2}\sum_{i<j} \ell^{\,i\frac{\vec{k}\cdot\vec{R}A}{\sqrt{2}}} \sum_{+,-} \pm\left(\ell^{\,i\frac{\vec{k}\cdot\vec{r}A}{\sqrt{2}}} \pm \ell^{\,-i\frac{\vec{k}\cdot\vec{r}A}{\sqrt{2}}}\right)(\vec{\sigma}_i \pm \vec{\sigma}_j)\cdot\hat{r}_{ij}$$
$$\{C_1 Y_1(Am_\pi r) + C_2 F_0(r)\}\, i(\vec{\tau}_i \times \vec{\tau}_j)_-, \qquad (A.1)$$

$$\vec{R} = \frac{1}{b\cdot\sqrt{2}}(\vec{r}_i + \vec{r}_j).$$

The antisymmetrized wavefunction of two particles in the well of the spherical harmonic oscillator is transformed from the absolute coordinates to the relative and centre-of-mass coordinates according to the prescriptions of the classical shell model by use of the Brody-Moshinsky coefficients [20]. We decompose the operator (A.1) into spherical tensors and apply the standard Racah-algebra technique in order to obtain the reduced matrix elements:

$$(n\ell S(j);\, NL;\, JT \| J^{4(\mu,\beta^*)}_{\mathrm{exch}} \| n'\ell' S'(j');\, N'L';\, JT') =$$

$$18\sqrt{2}\sum_{\substack{i;\\ g=1,2}} (-)^{D_i} \begin{pmatrix} D_i & 1 & A_i \\ 0 & 0 & 0 \end{pmatrix} (-)^{J+j'+\ell+S'} \xi_i [1+\xi_i(-)^{S+S'}]\, \hat{J}\hat{A}_i\hat{A}_i\hat{D}_i\hat{D}_i$$

$$\hat{S}\hat{S}'\hat{T}\hat{T}'\hat{j}'\hat{j}\hat{\ell}\hat{\ell}'\hat{L}\hat{L}' \begin{pmatrix} \ell & A_i & \ell' \\ 0 & 0 & 0 \end{pmatrix} \begin{pmatrix} L & D_i & L' \\ 0 & 0 & 0 \end{pmatrix} \begin{Bmatrix} j & j' & D_i \\ L' & L & J \end{Bmatrix} \begin{Bmatrix} 1 & S & S' \\ 1/2 & 1/2 & 1/2 \end{Bmatrix} \qquad (A.2)$$

$$\begin{Bmatrix} T & T' & 1 \\ 1/2 & 1/2 & 1 \\ 1/2 & 1/2 & 1 \end{Bmatrix} \begin{Bmatrix} A_i & 1 & D_i \\ \ell & S & j \\ \ell' & S' & j' \end{Bmatrix} C_g \int_0^\infty R_{n\ell}(r)\, j_{D_i}(Q^{(\mu,\beta)} r)\, \phi_g(r) R_{n'\ell'}(r) \times r^2 dr$$

$$\int_0^\infty R_{NL}(r')\, j_{D_i}(Q^{(\mu,\beta)} r')\, R_{N'L'}(r') r'^2 dr'.$$

In practical calculations we keep only $i = 1,2$.

$$Q^{(\mu,\beta)} = \frac{k^{(\mu,\beta)}}{\sqrt{2}}, \quad \begin{matrix} i = 1;\ A_i = 1,\ D_i = 0,\ \xi_i = 1 \\ i = 2;\ A_i = 0,\ D_i = 1,\ \xi_i = -1 \end{matrix}, \quad \phi_g(r) = \begin{cases} g = 1 & Y_1(Am_\pi r) \\ g = 2 & F_0(r) \end{cases}.$$

Here $R_{ab}(r)$ denotes the normalized radial function of the harmonic oscillator, $j_L(x)$ is the regular spherical Bessel function and $\hat{L} \equiv \sqrt{2L+1}$; $\begin{pmatrix} j_1 & j_2 & j_3 \\ m_1 & m_2 & m_3 \end{pmatrix}$, $\begin{Bmatrix} j_1 & j_2 & j_3 \\ \ell_1 & \ell_2 & \ell_3 \end{Bmatrix}$, and $\begin{Bmatrix} j_{11} & j_{12} & j_{13} \\ j_{21} & j_{22} & j_{23} \\ j_{31} & j_{32} & j_{33} \end{Bmatrix}$ are the 3j, 6j and 9j Wigner's symbols. For further notations see ref.[20].

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Мезонные обменные поправки к плотности ядерного слабого аксиального заряда в модели жестких пионов и $0^+ \leftrightarrow 0^-$ переход в ядрах с A = 16

Двухчастичная часть временной компоненты слабого аксиально-векторного тока строится в модели жестких пионов с минимальным кирально-инвариантным феноменологическим лагранжианом в приближении деревьев. Рассматриваются графики с обменом пиона, ρ и A_1 мезона. Полученный обменный оператор применяется для описания слабого чисто аксиального перехода $0^+ \leftrightarrow 0^-$, $\Delta T = 1$ в ядрах с A = 16. Корреляционные эффекты ядерной структуры описываются при помощи волновых функций модели оболочек со смешиванием конфигураций. Установлено сильное возрастание плотности ядерного слабого аксиального заряда по сравнению с импульсным приближением.

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Meson Exchange Corrections to Nuclear Weak Axial Charge Density in Hard Pion Model and $0^+ \leftrightarrow 0^-$ Transition in A = 16 Nuclei

Starting with the hard pion model based on a minimal chiral invariant phenomenological Lagrangian, the two-body part of the time component of the weak axial-vector current is constructed in the tree-approximation. Pion, rho- and A_1-meson exchanges are considered. The mesonic exchange operator obtained is applied to describe the purely weak axial $0^+ \leftrightarrow 0^-$, $\Delta T = 1$ transition in the nuclear A = 16 system. In order to treat nuclear structure correlation effects explicit use of shell model wave functions with configuration mixing is made. We confirm the large enhancement of the nuclear weak axial charge density with respect to impulse approximation.

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

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