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PIONIC DOUBLE CHARGE EXCHANGE WITHIN QUASIPARTICLE RANDOM PHASE APPROXIMATION

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The description of the pionic double charge exchange reaction is given within the proton-neutron quasiparticle random approximation. The approach is tested in a case of iron ^{56}Fe, and a fairly good agreement of the calculated quantities with recent data is found. The observed resonance-like behaviour of the energy dependence of the cross section is explained semiquantitatively in terms of two-nucleon process without invoking exotic mechanisms, like dibayrons or multiple quark clusters.

Дано описание реакции двойной перезарядки пионов в рамках протон-нейтронного квазичастичного приближения случайных фаз. Метод апробирован на примере ядра железа ^{56}Fe, и получено довольно хорошее согласие рассчитанных характеристик с современными данными. Наблюдаемое резонансноподобное поведение энергетической зависимости сечения полукачественно объяснено с помощью двухнуклонных процессов без привлечения экзотических механизмов, таких как дибарионы или многокварковые кластеры.

1. INTRODUCTION

Investigations concerning double charge exchange (DCX), both experimental and theoretical ones, have attracted much interest during the last decade. Studies of double charge exchange of pions on nuclei are attractive for many reasons. Since the charge conservation law ensures that at least two nucleons must be involved in pionic DCX on a nucleus, the reaction can be regarded as a promising source of specific information about the short range correlations between bound nucleons. Some authors also hope that the reaction makes it possible to study the expected difference between the neutron and proton densities in nuclei. Depending on the choice of the target nucleus, the DCX process may populate neutron- or proton-rich nuclei far from the stability region [1—4]. It is also possible to obtain information from the reaction about double isobaric analogue states [5,6], double isovector dipole resonances [7,8] or some exotic states of nuclear matter such as states of three and four neutrons [9,10]. (For a recent competent review of the pionic charge exchange reactions see Ref.11.)

In most theoretical investigations of the DCX reaction in the low energy region one assumes that the process is sequential: If two nucleons are correlated in space, then one can expect that a neutral pion π^0 emitted from the first charge exchange reaction on one nucleon finds a good chance to initiate the second charge changing process on the correlated partner. Further, one expects that the pion interacts during the two step exchange only with valence nucleons and thus the core plays only a passive role. This picture is quite natural for the transitions to the final double analogue state since during such a transition all quantum numbers of nucleons are unchanged except the third component of isospin (change of a neutron into a proton). In a non-analogue transition like the transition to the ground state of $T \neq 1$ nuclei, the core can play an active role due to the antisymmetrization of the total wave function, which allows the core nucleons to participate actively in the reaction [12].

Most of the theoretical approaches in the low pion energy domain are plane wave impulse approximation — PWIA theories. Some of them account for distortion effects (distorted wave impulse approximation — DWIA, coupled channel techniques). The conclusions are not unique and we can find statements about either importance [15—17] or negligence of the distortion [13—14]. Moreover, all calculations of DCX differential cross sections and angular distributions taking into account even simple correlated nuclear wave functions show a fairly satisfactory agreement with data, while theories without such correlations disagree with experiments by an order of magnitude or more. Some authors argue [13,14,18—20], that correlation effects are so important that it is impossible to see other effects such as those of reaction dynamics or the pion distortion unless nucleon-nucleon short range correlations will be accounted properly. It is also not clear what roles play the initial and final state interactions in this context.

From this point of view the nuclear structure involved in the DCX models is a very sensitive aspect of theoretical interpretations of the process. Here one can find pronounced differences ranging from very simple shell model approaches [13,14], through the generalized seniority scheme [21], to very advances *realistic* treatments [22,23]. Of course, the problem of nuclear structure will display its complexity as one deals with heavier nuclei. This is the reason why most existing DCX theoretical treatments concentrate on light nuclei. Recently, Vergados [24] proposed a treatment of the DCX reactions in the context of any shell model in which one separates reaction amplitudes into two parts, one depends on the nuclear wave function and another one is connected with characteristics of the charge changing process between a pion and a nucleon. It can help in some aspects, but even here one has the problem of explicit construction of the excited states of the intermediate nucleus and the Green functions.

In a series of papers [22,23,25] a new approach to the DCX process in the framework of the proton-neutron quasiparticle random phase approximation (pn-QRPA) was developed. The model utilizes wave functions in large configuration spaces for both protons and neutrons in the initial, intermediate and final nuclei and there is no need for a closure approximation.

The study of the DCX process is especially interesting for the medium-heavy nuclei, for which numerous data exist and therefore this case provides one more subtle test of the theory. Another point is the observation of the *resonance-like* behaviour around the pion energy $T_\pi = 50$ MeV. In contrast to the experimental observations microscopic treatments did not predict a rise of cross sections in this energy domain. Only recently Schepkin proposed a non-nucleon (dibaryon) mechanism [26,27] as a partial explanation of such a behaviour of low-energy cross sections. I am going to show that there is a chance to shed some light on this intriguing behaviour in the framework of an ordinary two-nucleon mechanism.

2. DOUBLE CHARGE EXCHANGE PROCESS

2.1. Brief Description of Chosen Charge Exchange Operators. The formalism of the charge exchange process for the low energy pions was applied to the DCX reactions on calcium [22], germanium [25] and tellurium [23] targets. I shall brifly recapitulate the main features of the theory, however in a more general form than in previous papers.

Within the simplest local πNN interaction Lagrangian known as pseudoscalar coupling, one can construct in the nonrelativistic approximation the effective pion-nucleon Hamiltonian [23,29]

$$h_p(\mathbf{q}) = -\sqrt{2}\, i \frac{f}{m_\pi} \Psi_N^\dagger \boldsymbol{\sigma} \cdot \mathbf{q}\, e^{i\mathbf{q}\cdot\mathbf{x}}\, \tau_+ \Psi_N , \qquad (1)$$

which generates the p-wave pion-nucleon interaction only. In eq. (1) $\boldsymbol{\sigma}$ and τ_+ are Pauli and isospin raising operators, respectively. $\Psi_N^\dagger$ (Ψ_N) are nucleon creation (annihilation) field operators and the momentum transfer is taken to be $\mathbf{q}$.

A possible *s*-wave contribution to the πNN interaction is obtained from phenomenological considerations taking into account a composite meson exchange mechanism [23,29,30]. It has the form

$$h_s(\mathbf{q}) = 4\pi \frac{\lambda_1}{m_\pi} \sqrt{2}\, \Psi^\dagger_N \omega_q \tau_+ \Psi_N . \tag{2}$$

In both Hamiltonians the plane wave approximation for pion wave functions was used. The constants f and λ_1 are determined to reproduce the experimental data for nucleon-nucleon and nucleon-pion elastic scattering [29,30]. Of course, in the case of bound nucleons they can, in general, be modified and fitted separately in each DCX reaction of interest. For the purpose of preserving the important features of the model and because of the approximations used in the construction of the charge exchange operators, we will apply these constants with experimentally determined values $f/4\pi = 0.08$ and $\lambda_1 = 0.046$ (Ref. 30 and refs. cited there).

The second quantization procedure applied to the nucleon field allows one to express the interaction Hamiltonians in terms of creation and annihilation operators for the protons ($c^\dagger_p$, c_p) and neutrons ($c^\dagger_n$, c_n) as:

$$h_p(\mathbf{q}) = -\sqrt{2}\, i \frac{f}{m_\pi} \sum_{pn} \left[\int d^3x \psi^*_p(\mathbf{x}) \sigma\cdot\mathbf{q}\, e^{i\mathbf{q}\cdot x}\, \psi(\mathbf{x}) \right] c^\dagger_p c_n \tag{3}$$

and

$$h_s(\mathbf{q}) = 4\pi \frac{\lambda_1}{m_\pi} \sqrt{2}\, \omega_q \sum_{pn} \left[\int d^3x \psi_p(\mathbf{x})^* \psi(\mathbf{x}) \right] c^\dagger_p c_n . \tag{4}$$

Here $\psi_a(\mathbf{x})$ is the solution of Schrödinger equation for any average nuclear potential, e.g., harmonic oscillator or Woods — Saxon with $a = p$ or n for protons and neutrons, respectively.

2.2 Transformation to Quasiparticles. Because of the quasiparticle character of the RPA, which we will use to describe the structure of the nuclei involved in the charge changing process one needs to transform expressions (3) and (4) to the Bogoliubov — Valatin quasiprotons ($a^\dagger_p$,a_p) and quasineutrons ($b^\dagger_n$, b_n)

$$a^\dagger_{pm_p} = u_p c^\dagger_{pm_p} + v_p (-1)^{j_p + m_p} c_{p-m_p} , \quad a_{pm_p} = (a^\dagger_{pm_p})^\dagger , \tag{5}$$

$$b^\dagger_{nm_n} = u_n c^\dagger_{nm_n} + v^p (-1)^{j_n + m_n} c_{n-m_m} , \quad b_{nm_n} = (b^\dagger_{nm_n})^\dagger , \tag{6}$$

u and v coefficients are related in a well-known way $u^2_{p(n)} + v^2_{p(n)} = 1$. After transformations (5)—(6) are performed, the p- and s-wave Hamiltonians have the form:

$$h_p(\mathbf{q}) = -\sqrt{2}\, i \frac{f}{m_\pi} \sum_{pn,\, JM} \mathcal{F}^{JM}_{pn}(\mathbf{q})\, \mathcal{R}^{JM}_{pn}\,, \tag{7}$$

$$h_s(\mathbf{q}) = -4\pi \frac{\lambda_1}{m_\pi} \sqrt{2}\, \omega_q \sum_{pn} (2j_p + 1)^{1/2}\, \delta_{pn}\, \mathcal{R}^{00}_{pn}\,. \tag{8}$$

In both eqs. (7) and (8) the transition density operator $\mathcal{R}^{JM}_{pn}$ is given by a formula

$$\mathcal{R}^{JM}_{pn} = u_p v_n C^\dagger(pnJM) + v_p u_n \tilde{C}(pnJM) + u_p u_n D(pnJM) - v_p v_n \tilde{D}^\dagger(pnJM). \tag{9}$$

The proton-neutron pair creation and annihilation operators and additional one-body type operators needed in the construction of operator (9) are defined in the usual way,

$$C^\dagger(pnJM) = \sum_{m_p,\, m_n} (j_p m_p j_n m_n \mid JM)\, a^\dagger_{pm_p} b^\dagger_{nm_n}\,, \tag{10}$$

$$C(pnJM) = [C^\dagger(pnJM)]^\dagger, \quad \tilde{C}(pnJM) = (-1)^{J+M}\, C(pnJ-M), \tag{11}$$

$$D^\dagger(pnJM) = \sum_{m_p m_n} (j_p m_p j_n m_n \mid JM) a^\dagger_{pm_p} (-1)^{j_n + m_n} b_{nm_n}\,, \tag{12}$$

$$D(pnJM) = [D^\dagger(pnJM)]^\dagger, \quad \tilde{D}(pnJM) = (-1)^{J+M} D(pnJ-M). \tag{13}$$

In eq. (7) we follow Ref. 23 for a definition of the function $\mathcal{F}^{JM}_{pn}$,

$$\mathcal{F}^{JM}_{pn}(\mathbf{q}) = \sum_{(m)} (-1)^{j_n - m_n} (j_p m_p j_n - m_n \mid JM) \left[\int d^3x \psi^*_p(\mathbf{x})\, \boldsymbol{\sigma}\cdot\mathbf{q} e^{i\mathbf{q}\cdot\mathbf{x}} \psi_n(\mathbf{x}) \right]. \tag{14}$$

After some algebra the function $\mathcal{F}^{JM}_{pn}$ can be written in a more compact form

$$\mathcal{F}^{JM}_{pn}(\mathbf{q}) = \sqrt{4\pi}\, \sqrt{6}\, Y^*_{JM}(\Omega_q) G^J_{pn}(q) \tag{15}$$

by setting apart the form factor G^J_{pn}:

$$G^J_{pn}(q) = (-1)^{j_p + j_n} \hat{j}_p \hat{l}_p \hat{j}_p \hat{l}_n \sum_{l''=0} (-1)^{\frac{l_p + l_n - l''}{2}} (l_p 0 l_n 0 \mid l'' 0)(J010 \mid l''0) \times \times R^{l''}_{pn}(q) \begin{Bmatrix} \frac{1}{2} & l_p & j_p \\ \frac{1}{2} & l_n & j_n \\ 1 & l'' & J \end{Bmatrix}. \tag{16}$$

In the above equation the symbol { } is the 9j-symbol (Fano) and (1) represents a Clebsh — Gordan coefficient. Further, Y_{JM} denotes the spherical harmonic depending on the solid angle Ω_q. The coefficient $\sqrt{6}$ in eq. (15) comes from the reduced matrix element of the nuclear spin operator $\boldsymbol{\sigma}$. We also used the abbreviation $\hat{j} = \sqrt{2j+1}$ and the corresponding expression for $\hat{l}$. In eq. (16) we introduced the overlap integral $R^{l''}_{pn}$ between radial nuclear wave functions and the radial part of the plane wave pion,

$$R^{l''}_{pn}(q) - \int_0^\infty dr\, dr^2 j_{l''}(qr)\, R_{n_p l_p}(r)\, R_{n_n l_n}(r). \tag{17}$$

Its explicit form depends on the choice of radial nucleon wave functions $R_{n_p l_p}$ and $R_{n_n l_n}$. In the case of harmonic-oscillator wave functions one can find a very compact analytical expression for this integral [23].

3. QRPA MODEL FOR THE DCX REACTION

3.1. DCX Amplitude and Intermediate Excited States. In second order perturbation theory the DCX transition amplitude is given as [23,25,37]

$$F_{if}(\mathbf{k}, \mathbf{k}') =$$

$$= \sum_{mm'} \frac{\langle f, 0^+; \pi^-(\mathbf{k}') \mid \hat{O} \mid mJM \rangle \langle mJM \mid \overline{m'JM} \rangle^* \langle \overline{m'JM} \mid \hat{O} \mid i(\text{G.S.}); \pi^+(\mathbf{k}) \rangle}{D(E_i, E^J_m, E^J_{m'}, \mathbf{q})}. \tag{18}$$

In the above equation $\mid i(\text{G.S.}), \pi^+(\mathbf{k})\rangle$ denotes the ground state of the initial nucleus (A, Z) and an incoming positive pion with momentum $\mathbf{k}$ and the initial energy $(k^2 + m_\pi^2)^{1/2}$. In analogy, $|f, \pi^-(\mathbf{k}')\rangle$ stands for an arbitrary state (ground or excited) of the final nucleus $(A, Z+2)$ and an outgoing negative pion with momenta $\mathbf{k}'$·$\mathbf{q}$ means the momentum transfer. Note that in expression (18) we have assumed that the charge operators are nonrelativistic Hamiltonians (3) and (4). It should be stressed that the denominator in eq. (18) differs in each case of the interaction (3) or (4). The Hamiltonian (3) represents a contribution of pion absorption on nuclear pair. From the general rules the Hamiltonian is known to be a small part of pair absorption at low energies. But in the DCX reaction it can play more important role [55]. The denominator in this case has a simple form $E_i + \omega_k - (E^J_m + E^J_{m'})/2$. The double scattering of a pion by two nucleons within sequential mechanism is caused by the interaction (4). Of course, the denominator for this channel has no explicit form and one can calculate the nuclear matrix elements of the charge changing operator including the propagator of the intermediate neutral pion. Integration over intermediate pion momentum is understood.

The DCX amplitude (18) contains all terms coming from two different sets of the intermediate states $\{|mJM\rangle\}$ and $\{|\overline{m'JM}\rangle\}$ generated within the proton-neutron quasiparticle random phase approximation (pn — QRPA). In principle, the two sets obtained by QRPA on the initial and on the final nucleus are identical and describe the excited states of the same intermediate nucleus $(A, Z+1)$. But, since both calculations are not exact solutions of the many-body problem, they lead to slightly different solutions $\{|mJM\rangle\}$ and $\{|\overline{m'JM}\rangle\}$ for the intermediate wave functions.

$$|mJM\rangle = Q^{m\,\dagger}_{JM}\,|RPA;(A,Z)\rangle \equiv$$

$$\equiv \sum_{(pn)}\left[X^{m}_{(pn)J}\,C^{\dagger}(pnJM) - Y^{m}_{(pn)J}\,\tilde{C}(pnJM)\right]|i;(\mathrm{G.S.})\rangle, \tag{19}$$

$$|\overline{m'JM}\rangle = \overline{Q}^{m'\,\dagger}_{JM}\,|RPA;(A,Z+2)\rangle \equiv$$

$$\equiv \sum_{(pn)}\left[\overline{X}^{m'}_{(pn)J}\,C^{\dagger}(pnJM) - \overline{Y}^{m'}_{pn)J}\,\tilde{C}(pnJM)\right]|f;(\mathrm{G.S.})\rangle. \tag{20}$$

Here $X(\overline{X})$ and $Y(\overline{Y})$ are forward- and backward-going amplitudes, respectively. p and n stand for proton and neutron quasiparticle states (compare eqs. (5) and (6)). As we pointed out states (19) and (20) are mathematically nonequivalent. In particular, the intermediate states belonging to different sets are not orthogonal and this is a reason for involving their overlaps $\langle mJ_m|\overline{m'J_{m'}}\rangle$ into sum (18). Some authors applied a similar procedure in double beta decay calculations [35,36]. We also used this scheme in our previous description of the DCX processes on calcium, germanium and tellurium isotopes [22,23,25].

In eq. (18) E_i is the initial energy of the parent (target) nucleus. Usually one can also adopt the average QRPA excitation energies $(E^J_m + E^J_{m'})/2$ in the denominator according to the above-mentioned procedure of accounting for the nonequivalence of the two sets of intermediate states. It is worth emphasizing the fact that the amplitude (18) does not contain the usual closure approximation in which one takes, instead of a sum over states in the odd-odd mass nucleus with their individual energies, some average energy equivalent for all states. In our calculations we use explicitly the intermediate QRPA states, their structure and the corresponding excitation energies in the denominator of eq. (18). In this respect we are able to find individual contributions coming from the different multipolarities J^{π} and to estimate the importance of the analogue and nonanalogue routes in the DCX reaction.

3.2. Excited States of the Daughter Nucleus. In Refs. 22 and 23 one can find expressions for the DCX amplitude in the case of ground and analogue state transitions. There are no principal difficulties to obtain more general formulae for transitions to any state of the final $(A, Z+2)$ nucleus. Analogously to what was done above (eqs. (19) and (20)), one can generate such states by the following *Anzatz*:

$$|\nu\mathcal{JM}\rangle = \overline{Q}^{\nu\dagger}_{\mathcal{JM}} |RPA; A+2)\rangle \equiv$$

$$\equiv \left\{ \sum_{(pp')} \left[\overline{\mathcal{X}}^{\nu}_{pp')\mathcal{J}} A^{\dagger}(pp'\mathcal{JM}) - \overline{\mathcal{Y}}^{\nu}_{(pp')\mathcal{J}} \tilde{A}(pp')\mathcal{JM}) \right] + \right.$$

$$\left. + \sum_{(nn')} \left[\overline{\mathcal{X}}^{\nu}_{(nn')\mathcal{J}} B^{\dagger}(nn'\mathcal{JM}) - \overline{\mathcal{Y}}^{\nu}_{(nn')\mathcal{J}} \tilde{B}(nn'\mathcal{JM}) \right] \right\} |f; \text{G.S.}\rangle. \qquad (21)$$

The creation and annihilation pair operators for protons $A^{\dagger}$, A and neutrons $B^{\dagger}$, B are defined in full analogy with eq. (10). $\overline{\mathcal{X}}^{\nu}$ and $\overline{\mathcal{Y}}^{\nu}$ stand for forward- and backward-going amplitudes. They, as well as other amplitudes, X, Y, $\overline{X}$, $\overline{Y}$, are determined by solving the appropriate QRPA equation of motion for the states in the initial, the intermediate and the final nucleus. Details of the structure of the QRPA equations and their solutions can be found in Refs. 22 and 23.

3.3. Matrix Elements and DCX Cross Section. Using the above expressions for Hamiltonians (3) and (4) and definitions of the intermediate (eqs. (19)—(20)) and final (eq. (21)) states, one can find the following formulae for matrix elements needed to write amplitude (18) in an explicit form:

— matrix element for the s-wave charge changing operator contributing to the transition between the intermediate state $|m'J^{\pi}M\rangle$ in the nucleus $(A, Z+1)$ and the final state $|\nu\mathcal{JM}\rangle$ of the daughter nucleus $(A, A+2)$

$$\langle \nu\mathcal{JM}; \pi^{-}(\mathbf{k}') | h_s | \overline{m'J^{\pi}M} \rangle = 4\pi \frac{\lambda_s}{m_{\pi}} \sqrt{2}\, \omega_k{}' \,\delta_{J\mathcal{J}} \delta_{M\mathcal{M}} \times$$

$$\times \sum_{p \le p', n''} (-1)^{j_p + j_{n''}} \left[\left(\overline{\mathcal{X}}^{\nu}_{(pp')\mathcal{J}} \overline{X}^{m'}_{(p'n'')J} \overline{u}_p \overline{u}_{n''} - \overline{\mathcal{Y}}^{\nu}_{(pp')\mathcal{J}} \overline{Y}^{m'}_{(p'n'')J} \overline{v}_p \overline{v}_{n''} \right) \delta_{pn''} + \right.$$

$$\left. + (-1)^{J} \left(\overline{\mathcal{X}}^{\nu}_{(pp')\mathcal{J}} \overline{X}^{m'}_{(pn'')J} \overline{u}_{p'} \overline{u}_{n''} - \overline{\mathcal{Y}}^{\nu}_{(pp')\mathcal{J}} \overline{Y}^{m'}_{(pn'')J} \overline{v}_{p'} \overline{v}_{n''} \right) \delta_{p'n''} \right]; \qquad (22)$$

— matrix element for the p-wave charge changing operator contributing to the transition between the intermediate state $|\overline{m'J^{\pi}M}\rangle$ and the final state $|\nu\mathcal{JM}\rangle$

$$\langle \nu\mathcal{JM}; \pi^-(\mathbf{k}') \,|\, h_p \,|\, \overline{m'J^{\pi}M} \rangle =$$

$$= i\frac{f}{m_\pi}\sqrt{4\pi}\sqrt{12}\sum_{p\le p',n''J''M''} \hat{J}\hat{J}''(JMJ''M''\,|\,\mathcal{JM})\,(-1)^{j_p+j_{p''}+J+J''}\,Y^*_{J''}(\Omega_{k'})\times$$

$$\times\left[\begin{Bmatrix} j_{n''} & j_{p'} & J \\ \mathcal{J} & J'' & j_p \end{Bmatrix} G^{J''}_{pn''}(k')\left(\overline{X}^{\nu}_{(pp')\mathcal{J}}\overline{X}^{m'}_{(p'n'')J}\bar{u}_p\,\bar{u}_{n''} - \overline{\mathcal{Y}}^{\nu}_{(pp')\mathcal{J}}\overline{Y}^{m'}_{(p'n'')J}\bar{v}_p\,\bar{v}_{n''}\right)+\right.$$

$$\left.+(-1)^{\mathcal{J}}\begin{Bmatrix} j_{n''} & j_p & J \\ \mathcal{J} & J'' & j_{p'} \end{Bmatrix} G^{J''}_{p'n''}(k')\left(\overline{X}^{\nu}_{(pp')\mathcal{J}}\overline{X}^{m'}_{(pn'')J}\bar{u}_{p'}\,\bar{u}_{n''} - \overline{\mathcal{Y}}^{\nu}_{(pp')\mathcal{J}}\overline{Y}^{m'}_{(pn'')J}\bar{v}_{p'}\,\bar{v}_{n''}\right)\right]. \quad (23)$$

Formulae (22)—(23) are written only for proton-proton quasiparticle excitations in the final nucleus. In the complete expressions one has to add analogous terms for neutron-neutron excitations.

Further, we also need two other matrix elements:

— matrix element for the *s*-wave charge changing operator contributing to the transition between the ground state $|i;(\mathrm{G.S.})\rangle$ of the parent nucleus and the intermediate states $|mJ^{\pi}M\rangle$

$$\langle mJ^{\pi}M \,|\, h_s \,|\, i;(\mathrm{G.S.});\pi^+(\mathbf{k})\rangle =$$

$$= \sqrt{2}\left(4\pi\frac{\lambda_s}{m_\pi}\right)\omega_k\delta_{J0}\delta_{M0}\left[\sum_{pn}\delta_{pn}\left(X^m_{(pn)J}u_p v_n + Y^m_{(pn)J}v_p u_n\right)\right]; \quad (24)$$

— matrix element for the *p*-wave charge changing operator contributing to the transition between the initial ground state $|i;(\mathrm{G.S.})\rangle$ of the parent nucleus and the intermediate states $|mJ^{\pi}M\rangle$

$$\langle mJ^{\pi}M \,|\, h_{sp} \,|\, i;(\mathrm{G.S.});\pi^+(\mathbf{k})\rangle =$$

$$= i\sqrt{12}\frac{f}{m_\pi}\sqrt{4\pi}\,(-1)^J\,Y^*_{JM}(\Omega_k)\left[\sum_{pn}G^J_{pn}(k)\left(X^m_{(pn)J}u_p v_n + Y^m_{(pn)J}v_p u_n\right)\right]. \quad (25)$$

Expressions and definitions (18), (22)—(25) allow one to write down the explicit formula for the DCX amplitudes in the most general case of the transition to any final state $|\nu\mathcal{JM}\rangle$,

$$F_J^s(\mathbf{k},\mathbf{k}') = -\delta_{J\mathcal{J}}\left(4\pi\frac{\lambda_s}{m_\pi}\right)^2 \omega_k\,\omega_{k'} \sum_{m,m'} \frac{\langle mJ^\pi | \overline{m'J^\pi}\rangle}{D(E_i, E_m^J, E_{m'}^J, \mathbf{q})} \times$$

$$\times\Bigg\{\sqrt{2}\sum_{n'',p\le p'}(-1)^{j_p+j_{n''}}\Bigg[\left(\overline{\mathcal{X}}^{\nu}_{(pp')\mathcal{J}}\overline{X}^{m'}_{(p'n'')J}\,\bar{u}_p\bar{u}_{n''} - \overline{\mathcal{Y}}^{\nu}_{pp'\mathcal{J}}\overline{Y}^{m}_{(p'n''J}\,\bar{v}_p\bar{v}_{b''}\right)\delta_{pn''} +$$

$$+(-1)^J\left(\overline{\mathcal{X}}^{\nu}_{(pp')\mathcal{J}}\overline{X}^{m'}_{(pn'')J}\,\bar{u}_{p'}\bar{u}_{n''} - \overline{\mathcal{Y}}^{\nu}_{(pp')\mathcal{J}}\overline{Y}^{m}_{(pn'')}\,\bar{v}_{p'}\bar{v}_{n''}\right)\delta_{p'n''}\Bigg] +$$

$$+\sum_{p'',n\le n'}\left[\begin{pmatrix} p\to n\\ p'\to n'\\ n''\to p''\end{pmatrix}\right]\Bigg\}\left[\sqrt{2}\,\delta_{J0}\sum_{pn}\hat{j}\,\delta_{pn}\left(X^m_{(pn)J}u_p v_n - Y^m_{(pn)J}v_p u_n\right)\right]. \tag{26}$$

$$F_J^p(\mathbf{k},\mathbf{k}') = -\left(\frac{f}{m_\pi}\right)^2 \sum_{m,m'} \frac{\langle mJ^\pi | \overline{m'J^\pi}\rangle}{E_i + \omega_k - \dfrac{E_m^J + E_{m'}^J}{2}} \times$$

$$\times\sum_{J''\mathcal{M}} \hat{\mathcal{J}}\hat{J}''(-1)^{\mathcal{J}+\mathcal{M}}\{Y_{J''}(\omega_{k'})\otimes Y_{J''}(\Omega_k)\}_{\mathcal{J}\mathcal{M}}\Bigg\{\sqrt{12}\sum_{p\le p',n''}(-1)^{j_p+j_{p''}}\times$$

$$\times\Bigg[\begin{Bmatrix} j_{n''} & j_{p'} & J\\ \mathcal{J} & J'' & j_p\end{Bmatrix} G^{J''}_{pn''}(k')\left(\overline{\mathcal{X}}^{\nu}_{(pp')\mathcal{J}}\overline{X}^{m'}_{(p'n'')J}\bar{u}_p\,\bar{u}_{n''} - \overline{\mathcal{Y}}^{\nu}_{(pp')\mathcal{J}}\overline{Y}^{m'}_{(p'n'')J}\,\bar{v}_p\,\bar{v}_{n''}\right)+$$

$$+(-1)^{\mathcal{J}}\begin{Bmatrix} j_{n''} & j_p & J\\ \mathcal{J} & J'' & j_{p'}\end{Bmatrix} G^{J''}_{p'n''}(k')\left(\overline{\mathcal{X}}^{\nu}_{(pp')\mathcal{J}}\overline{X}^{m'}_{(pn'')J}\bar{u}_{p'}\,\bar{u}_{n''} - \overline{\mathcal{Y}}^{\nu}_{(pp')\mathcal{J}}\overline{Y}^{m'}_{(pn'')J}\,\bar{v}_{p'}\,\bar{v}_{n''}\right)\Bigg]+$$

$$+\sqrt{12}\sum_{p'',n\le n'}\left[\begin{pmatrix} p\to n\\ p'\to n'\\ n''\to p''\end{pmatrix}\right]\Bigg\}\left[\sqrt{12}\sum_{pn} G^J_{(pn)J}\left(X^{m'}_{(pn)J}u_p v_n - Y^{m'}_{(pn)J}v_p u_n\right)\right]. \tag{27}$$

The full amplitude $F_{\text{if}}(\mathbf{k},\mathbf{k}')$ is taken to be a sum of all multipolarities allowed for each part of the charge changing operator. Selection rule steming from angular momentum and parity conservation laws limits the transition caused by s-wave operator (4) only to 0^+ intermediate states, whereas p-wave operator (3) has nonzero contributions for all the so-called pion-like intermediate states (0^-, 1^+, 2^-, 3^+...). The differential cross section is normalized in such a way that

$$\frac{d\sigma}{d\Omega}(\theta, q) = \left|\frac{1}{4\pi}\sum_{J^\pi}[F_J^s(\mathbf{k},\mathbf{k}') + F_J^p(\mathbf{k},\mathbf{k}')]\right|^2, \tag{28}$$

where $\mathbf{q} = \mathbf{k} - \mathbf{k}'$ is the momentum transfer in the DCX process.

As is mentioned above we used throughout the paper the nonrelativistic approximation for the charge changing operator. In general, the relativistic corrections can influence the differential cross section (28). One can expect a negligible role of such terms for the p-wave part of the transition amplitude and a larger contribution to the s-part. The plane wave approximation for pions used in this paper can also be a source of inaccuracies for the cross section. I shall not treat the pion distortions in this paper.

3.4. Ground State DCX Transition. Formulae (26) and (27) represent general expression for the amplitude of the DCX transition into any final state of a daughter nucleus. The DCX transition to the ground state can be simplified to the amplitude [23,27]:

$$F^{s}_{\rm GS}(\mathbf{k},\mathbf{k}') = \left(4\pi\frac{\lambda_s}{m_\pi}\right)^2 \omega_k\,\omega_{k'} \sum_{mm'} \frac{\langle m0^+ \mid \overline{m'0^+}\rangle}{D(E_i, E^0_m, E^0_{m'}, \mathbf{q})} \times$$

$$\times\left\{\left[2\sqrt{2}\sum_{pn}\delta(p,n)\hat{j}_p\left(\overline{X}^{m'}_{(pn)0}\overline{v}_p\overline{u}_n - \overline{Y}^{m'}_{(pn)0}\overline{u}_p\overline{v}_n\right)\right]\times\right.$$

$$\left.\times\left[2\sqrt{2}\sum_{pn}\left(X^m_{(pn)0}u_pv_n - Y^m_{(pn)0}v_pu_n\right)\right]\right\}. \tag{29}$$

$$F^{p}_{\rm GS}(\mathbf{k},\mathbf{k}') = -\left(\frac{f}{m_\pi}\right)^2 \sum_{mm',J} \frac{\langle mJM \mid \overline{m'JM}\rangle}{E_i + \omega_k - \dfrac{E^J_m + E^J_{m'}}{2}} P_J(\cos\theta_{kk'}) \times$$

$$\times\left\{\left[\sqrt{12}\sum_{pn}(-1)^{j_p+j_n+J}G^J_{pn}(k')\left(\overline{X}^{m'}_{(pn)J}\overline{v}_p\overline{u}_n - \overline{Y}^{m'}_{(pn)J}\overline{u}_p\overline{v}_n\right)\right]\times\right.$$

$$\left.\times\left[\sqrt{12}\sum_{pn}G^J_{(pn)J}(k)\left(X^m_{(pn)J}u_pv_n - Y^m_{(pn)J}v_pu_n\right)\right]\right\}. \tag{30}$$

All symbols in the two last equations have been used already. The only new quantity $P_J(\cos\theta_{kk'})$ is Legendre polynomial coming from the reduction of the tensor product of spherical harmonics in eq. (27). We also used the abbreviation $\delta(p,n) = \delta_{n_p n_n}\delta_{j_p j_n}\delta_{l_p l_n}$.

We would like to note that the s-wave part of the transition operator contains only the route in which the neutron occupying some nuclear particle state with defined quantum numbers n, l, j is changed into the proton with exactly the same quantum numbers. These transitions excite the isobaric analogue state (IAS), if they are coherently superimposed. However, because of the pairing correlations the situation is more complicated and one obtains appreciable transitions through other 0^+ states. In the theory developed here we are able to separate both types of contributions using a method of identification of the IAS transition which, in general, represents the strongest 0^+-transition [23]. We already stressed that the p-wave part of the charge changing operator gives a contribution to the DCX amplitude only for the *pion-like* intermediate states and thus causes nonanalogue routes which are sensitive to short-range nucleon-nucleon correlations.

4. EXAMPLE OF THE DCX REACTION ON ^{56}Fe

4.1. Details of the Calculations. As an example of application of the theory I shall discuss the ground transition in iron: $^{25}\mathrm{Fe} \to {}^{56}\mathrm{Ni}$. This reaction was already studied in the eighties [38,39], but the data base is still very sparse for nonanalogue transitions. For the pion energy of 50 MeV three experimental points of the angular distribution have been measured at PSI for the ground to ground transition as well as to several individual predominantly 0^+ excited states of ^{56}Ni [38]. Some preliminary data also exist at pion energies $T_\pi = 35$ and 61 MeV. Although the data for the ground to ground transition is very limited, one already can conclude the following: The transition exhibits a well pronounced resonance-like behaviour. This means that the cross section at $T_\pi = 50$ MeV is for forward angles by one order of magnitude or more larger than at other energies.

The results presented below are achieved with a model space consisting of $0p_{1/2}$, $0p_{3/2}$, $1s_{1/2}$, $0d_{3/2}$, $0d_{5/2}$, $1p_{1/2}$, $1p_{3/2}$, $0f_{5/2}$ and $0f_{7/2}$. The single particle states used here are calculated with a Coulomb-corrected Woods — Saxon potential. It was assumed that both types of nucleons — protons and neutrons — occupy the same shells and calcium ^{40}Ca was taken as the inactive core. Two-body matrix elements needed for construction of the QRPA matrices were obtained from the realistic nuclear matter G-matrix by solving the Bethe — Goldstone equation (see, e.g., Refs. 40 and 41 for more details)

$$G(\omega) = V + \frac{Q}{\omega - H_0} G(\omega) . \tag{31}$$

In the above equation Q is the Pauli projection operator, ω stands for the starting energy and V is taken to be the nucleon-nucleon realistic one-meson exchange Bonn potential [42—44]. H_0 is the unperturbed single particle

Hamiltonian. In the present work we used the harmonic oscillator Hamiltonian. To take into account the effects of the finite nucleus we solve eq. (31) with as small absolute value of the starting energy as -25.0 MeV. This corresponds to an average single particle energy of -12.5 MeV. The oscillator length used is $b = 2.0$ fm.

The two-body matrix elements are obtained for nuclear matter. They are not specialized for a given nucleus. Thus and due to the finite Hilbert space used one has to renormalize them by multiplying with factors slightly different from 1.0: $g^n_{\rm pair}$, $g^p_{\rm pair}$, g^{pn}_{pp} and g^{pn}_{ph}. For the ground states of the parent and daughter nuclei one obtains uncorrelated vacuum states by solving the standard BCS equation in the above-mentioned model space. Two renormalization factors $g^n_{\rm pair}$ and $g^p_{\rm pair}$ multiplying the proton and neutron pairing matrix elements $<(aa)0|G|(bb)0>$ are fixed by adjusting the empirical pairing gaps Δ^p_a and Δ^n_a to the lowest quasiparticle energy obtained from the gap equation

$$\Delta^{p(n)}_a = \frac{1}{2}\, g^{p(n)}_{\rm pair}\, \hat{j}_a^{-1} \sum_b \hat{j}_b \Delta_b \left[(\varepsilon_b - \lambda_{p(n)})^2 + \Delta_b^2\right]^{-1/2} <(aa)0|G|(bb)0> . \qquad (32)$$

The empirical pairing gaps are deduced according to the recently published prescription of Moeller and Nix [45,46]:

$$\Delta^{\rm even-even}_{\rm neutron} = -\frac{1}{8}\,[M(Z,N+2) - 4M(Z,N+1) + 6M(Z,N) - \\ - 4M(Z,N-1) + M(Z,N-2)] , \qquad (33)$$

$$\Delta^{\rm even-even}_{\rm proton} = -\frac{1}{8}\,[M(Z+2,N) - 4M(Z+1,N) + 6M(Z,N) - \\ - 4M(Z-1,N) + M(Z-2,N)] . \qquad (34)$$

Expressions (33) and (34) cannot be used for nuclei with a magic number of protons or neutrons. Thus the pairing gap and the corresponding pairing strengths are estimated using the adjacent even-even nucleus. The table contains values of the pairing strengths $g^n_{\rm pair}$ and $g^p_{\rm pair}$ fitted in this way for ^{56}Fe and ^{56}Ni. All the renormalization factors are close to unity. Thus the bare G-matrix elements of the Bonn potential are already reasonably good.

Solutions of the BCS equations with matrix elements fixed in this manner allow one to evaluate the occupation amplitudes u's and v's needed for construction of the QRPA equation of motion. To determine the QRPA matrices fully one must also fix two additional renormalization factors, the strength of the particle-particle g^{pn}_{pp} and the strength of the particle-hole g^{pn}_{ph} interaction. For

Table. Experimental neutron and proton gaps for ^{56}Fe, ^{54}Fe and ^{58}Ni nuclei obtained from eqs. (33) and (34). The last two nuclei are used to estimate the strengths in nickel ^{56}Ni because of its double-magic character. (see text for details). Masses are taken from the mass tables [54]. The pairing strengths $g^n_{\rm pair}$ and $g^p_{\rm pair}$ were fixed to reproduce the experimental gaps

Nucleus	$\Delta^n_{\rm exp}$, MeV	$\Delta^p_{\rm exp}$, Mev	$g^n_{\rm pair}$	$g^p_{\rm pair}$
^{56}Fe	1.360	1.570	0.938	0.993
^{54}Fe	—	1.520	—	0.908
^{58}Ni	1.300	—	1.030	—

this purpose we used the isobaric state (IAS) and the Gamow-Teller state in cobalt ^{56}Co which are known to be 3.65 and 10.60 MeV, respectively [47,48]. The QRPA energy of these states depends predominantly on the particle-hole strength and adjustment of them to experimental energies gives $g^{pn}_{ph} = 0.8$. Details of such a procedure are given in Ref. 23. The second factor g^{pn}_{pp} will be treated as a free parameter of the theory and we will discuss all reaction observables as a function of it.

4.2. Results and Discussion. We calculated angular distributions and the energy dependence of the cross section for the DCX ground-ground transitions on ^{56}Fe. Figure 1 shows the angular distribution for the incident pion energy $T_\pi = 50$ MeV. Three curves are presented for three different values of the particle-particle strength g^{pn}_{pp}: 0.8, 1.0 and 1.1. The experimental points are measured at the Paul Scherer Institute by the Tübingen — Karlsruhe group [38]. The angular distribution decreases rapidly as the particle-particle parameter increases. This behaviour is observed in a full range of the g^{pn}_{pp} strength up to the value 1.1. for which the QRPA solution tends to a collapse. A similar behaviour was also observed in other nuclei [22,23,25] and other processes, e.g., the double-beta decay [35,36]. The mechanism for the collapse is connected with increase of the ground-state correlations by enlarging the particle-particle interaction. As a result, the lowest excited QRPA state is pushed down in energy below the ground state. Simultaneously the cross sections drop by factors of 3—10 depending on the scattering angle. The cross section is reduced since increasing g^{pn}_{pp} produces stronger the ground-state correlations. This enlarges the backward-going amplitudes Y's. The terms with Y's in eqs. (29), (30) become large enough to cancel against the terms with the forward-going

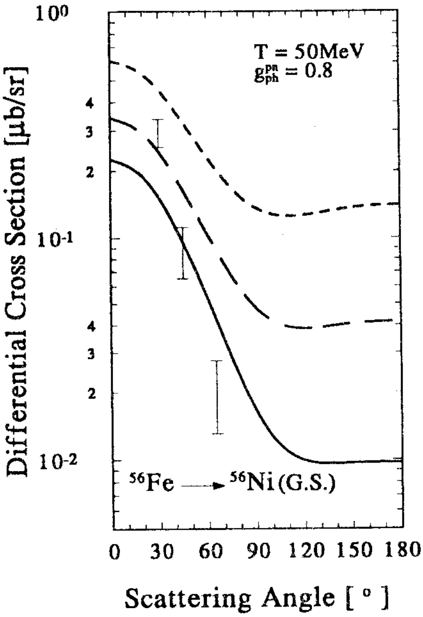

Fig. 1. Angular distribution for the ground state transition on iron ^{56}Fe. The particle-hole strength g_{ph}^{pn} is fixed to reproduce the Gamow — Teller and isobaric analogue state difference in ^{56}Ni. The results for three values of the particle-particle strength are shown: $g_{pp}^{pn} = 0.8$ (short-dashed line), $g_{pp}^{pn} = 1.0$ (long-dashed line) and $g_{pp}^{pn} = 1.1$ (soled line). The experimental data are taken from refs. 38 and 53

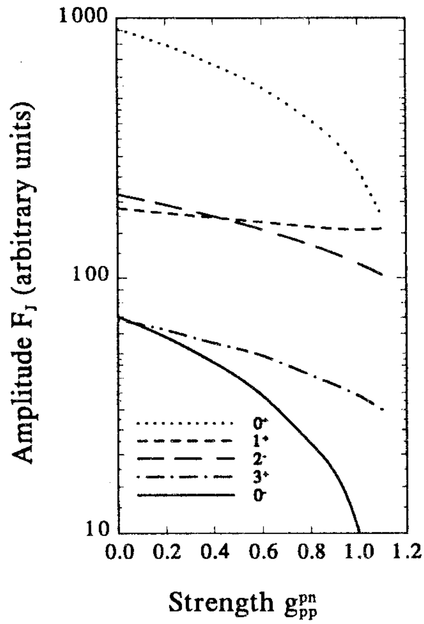

Fig. 2. The most important contributions to the transition amplitude come from the intermediate 0^+ (dotted line), 0^- (solid line), 1^+ (short-dashed line), 2^- (long-dashed line), and 3^+ (dashed-dotted line) states. The results for the pion energy $T_\pi = 50$ MeV are plotted as a function of the particle-particle strength g_{pp}^{pn}. The particle-hole strength is fixed to be 0.8

amplitudes X's. The magnitude of the DCX cross section rapidly diminishes. Comparison of the experimental results and theoretical predictions (fig.1) shows that the physically important domain of g_{pp}^{pn} is the interval 1.0—1.1.

It is interesting to compare the importance of contributions to the total amplitude coming from the different angular momenta. Such contributions for the forward-angle DCX amplitude and for the intermediate states $J^\pi = 0^+, 0^-, 1^+, 2^-$ and 3^+ which are most important, are presented in fig.2. One can notice immediately that a crucial role in pushing the contributions down

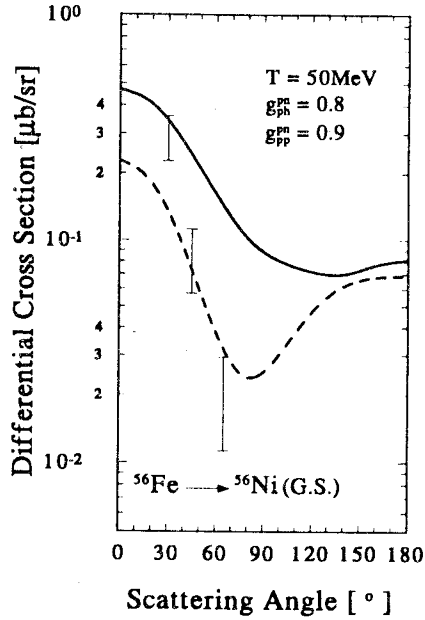

Fig. 3. Angular distributions for two choices of the model space are presented for fixed interaction strengths ($g_{ph}^{pn} = 0.8$, $g_{pp}^{pn} = 0.9$) and the incident pion energy $T_\pi = 50$ MeV. The solid line represents the results of a smaller and the dashed line of a very large model space. The experimental values are from Refs. 38 and 53. (For details see text.)

into agreement with the data is played by the dependence of each partial amplitude on the particle-particle strength g_{pp}^{pn}. The transition through 0^+ states with the biggest contribution from the IAS dominates the total amplitude at law g_{pp}^{pn}. But in the physically interesting interval $g_{pp}^{pn} \geq 1.0$ this amplitude is comparable with the 1^+ and 2^- amplitudes. Thus, all approaches restricted to the intermediate isobaric analogue state are not a good description of the cross sections and other DCX observables because nonanalogue routes play as an important role as transitions through 0^+ states.

We also examine in this paper the influence of the model space on the final results. Additional calculations of angular distributions at the pion energy $T_\pi =$ $= 50$ MeV were performed for a «huge» single particle basis consisting of the states $0s_{1/2}$, $0p_{1/2}$, $0p_{3/2}$, $0d_{3/2}$, $0d_{5/2}$, $1p_{1/2}$, $1p_{3/2}$, $0f_{5/2}$, $0f_{7/2}$, $2s_{1/2}$, $1d_{3/2}$, $1d_{5/2}$, $0g_{7/2}$, $0g_{9/2}$, $2p_{1/2}$, $2p_{3/2}$, $1f_{5/2}$, $1f_{7/2}$, $0h_{11/2}$, $3s_{1/2}$, $2d_{3/2}$, $2d_{5/2}$, $1g_{7/2}$, $1g_{9/2}$ for neutrons and $0s_{1/2}$, $0p_{1/2}$, $0p_{3/2}$, $1s_{1/2}$, $0d_{3/2}$, $0d_{5/2}$, $1p_{1/2}$, $1p_{3/2}$, $0f_{5/2}$, $0f_{7/2}$, $2s_{1/2}$, $1d_{3/2}$, $1d_{5/2}$, $0g_{7/2}$, $0g_{9/2}$, $2p_{1/2}$, $2p_{3/2}$, $1f_{7/2}$ for protons. All these single particle levels are below 5.0 MeV in Woods — Saxon potential and are either bound or quasi-bound. Figure 3 presents the angular distribution for this two choices of the size of the basis for the particle-particle strength $g_{pp}^{pn} = 0.9$. A change of the shape of the angular distribution by increasing the basis is clearly seen. A minimum around the scattering angle $\theta = 70°$ appears, which is also in the agreement with Gibbs' prediction [38]. Compared with the «small» basis, the angular distribution with the large basis is steeper. Absolute values of the

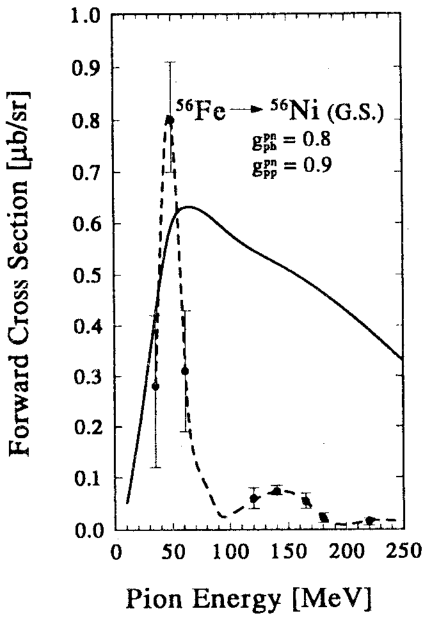

Fig. 4. Energy dependence of the forward-angle (5°) cross section for the ground state transition on iron ^{56}Fe (solid line). The results are obtained for the small basis. Data indicated by error bars are taken from Refs. 38 and 39 (the drawn short-dashed curve is only to guide the eye)

cross section are 2—5 times smaller for the same particle-particle strength g_{pp}^{pn}. This means one obtains agreement with the experiment for smaller values of the particle-particle strength. In this stage of the theory we are not able to separate two effects which influence the lowering of the cross section, i.e., particle-particle correlations manifested by the magnitude of g_{pp}^{pn} and participation of core nucleons in the DCX process (in the «huge» basis all nucleons are involved).

Recently measured ground-ground transition ^{56}Fe on at pion energies $T_\pi = 35$ and 61 MeV [38] together with the earlier data at higher energies from LAMPF [39] allow to systematize the dependence of the DCX cross sections. The experimental observations shown in fig.4 exhibit a resonance-like structure near $T_\pi \approx 50$ MeV.

In a contrast to such an observed behaviour almost all of the theoretical models with a distorted wave as well as with plane wave approximations are not able to predict even roughly this strong energy dependence for the DCX forward-angle cross section. The microscopic calculations give a rather smooth energy behaviour around $T_\pi = 50$ MeV except the predictions of Martemyanow and Schepkin [26,27] who have introduced dibayron resonance «by hand» to explain this dependence. These authors proposed a very narrow dibayron resonance formed in the DCX whose decay into two nucleons is not allowed by selection rules. The condition for building this dibaryon is a large overlap of a pair of nucleons (neutron-neutron or proton-proton) in their relative s-wave with $J^\pi = 0^+$ and $T = 1$. This resonance can appear according to refs. 26 and 27 at distances less than 1 fm between nucleons. Taking the estimation of G.Miller for the 6-quark bag probability to be of the order of a few per cent for all nucleon pairs in a nucleus [49], Martemyanow and Schepkin obtained an energy

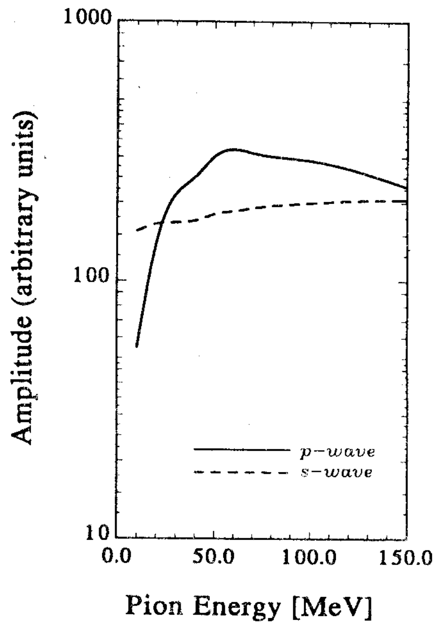

Fig. 5. Energy dependence of the analogue (dashed line) and non-analogue (full line) contributions to the ground state transition amplitude for the scattering angle $\theta = 5°$. The interactions strengths g^{pn}_{ph} and g^{pn}_{pp} are taken to be 0.8 and 1.1, respectively

dependence for the dibaryon mechanism in the double charge exchange reaction which roughly follows the observed behaviour. All such frameworks together with calculations made by Chiang and Zou [50] can be treated as an indication of the importance of quarks in the DCX process.

Presented approach, different in spirit, can also give the gross features of the energy dependence of the DCX. The calculated forward (5°) cross section is shown in fig.4 as a function of the pion energy up to the resonance region*. One can notice that we are able to explain qualitatively within our mechanism the observed experimental behaviour. The curve is not so steep on the high energy side above 60 MeV, but a peak around $T_\pi = 60$ MeV is clearly seen. It is rather obvious that we should look for some effects to reduce the DCX amplitude in the higher energy domain.

A more careful analysis of both *s*- and *p*-wave contributions to the total amplitude at two energies, say 10 and 50 MeV, and at forward angles may supply possible such a mechanism [51]. In fig.5 the dependence of the analogue and nonanalogue amplitude on the incident pion energy is shown. The *s*-wave contribution is almost constant in full domain of the pion energy. A dramatic increase of the nonanalogue amplitude (*p*-wave contribution) is seen up to the energy $T_\pi = 60$ MeV. So this component produces the maximum in the cross section. Because the nonanalogue route (*p*-wave component) depends sensitively on g^{pn}_{pp} [13,14,23] the strong dependence of the amplitude suggests

*Generalization of the model for the Δ-isobar degrees of freedom is possible [37], but we do not intend to discuss such a point in this paper.

that particle-particle correlations are of growing importance as one takes more energetic pions. After maximum is reached both components, s- and p-wave stay equally important. One clearly needs some additional mechanism for the reduction if the experimental data have to be reproduced. Such a possibility is offered by improving the s-wave charge changing operator by taking its relativistic form.

It is worth noting, that Karapiperis and Kobayashi [16] can roughly predict the decrease of the ^{14}C cross section between 50 and 100 MeV. Unfortunately cited calculations were not performed for pion energies lower than 50 MeV. Also Gibbs and co-workers proposed a resonance phenomenon in the pion scattering, which could be seen even more clearly in DCX around the proper energy $T_\pi = 50$ MeV [52]. These approaches — including presented here — point out a possibility of explaining the observed resonance-like behaviour without invoking nonstandard mechanisms. The existing data does not yet discriminate clearly between conventional and more exotic interpretations.

5. FINAL REMARKS

I have investigated the double charge exchange reaction in the framework of the quasi-particle random phase approximation (QRPA). The charge exchange operator was taken in the nonrelativistic form and the plane wave approximation was used for the incident, intermediate and outgoing pions.The approach was applied to the ground state transition on iron ^{56}Fe. The predicted values underestimate the forward-angle cross section and thus the calculated angular distribution is flatter than in the experiment.

Amplitudes and cross sections show smooth dependence on the value of the particle-particle strength g_{pp}^{pn}, which was also observed in earlier calculations of the DCX reaction on calcium [22], germanium [25] and tellurium [23]. A comparison with data allows one to state the physically important values of the strength g_{pp}^{pn} for iron and nickel nuclei lay in the interval 1.0—1.1. One should stress that the choice of the model space influences the calculated quantities. Because of this effect the g_{pp}^{pn} strength is not unique. As larger the basis as smaller is the particle-particle strength. Moreover, in the larger model space we observed a collapse of the QRPA solution for the g_{pp}^{pn} value as small as 0.9 which may suggest a need for inclusion of higher RPA-corrections into the model.

The gross features of the resonance-like shape of the cross section as a function of the pion energy can be reproduced at least semiquantitatively within the conventional $2N$ mechanism. The prediction is not too good, probably due to approximations made. It has to be seen in the future if the very speculative idea of a dibaryon resonance in the DCX reaction will prevail. Future development of the approach will make it possible to settle more carefully the questions addressed in this talk.

The last but not the least, a sensitivity of the pionic DCX processes to nuclear structure and especially to nucleon-nucleon correlations makes them interesting for the double beta decay. In searches for physics beyond the standard model the last reaction has continually received much attention. Grand unified theories predict the neutrinoless double beta decay if the neutrino is a Majorana particle with rest mass and/or the week right-handed currents exist. Combining both phenomena (DCX and double beta decay) ensures reliable nuclear matrix elements and thus an accurately defined estimate of the nonstandard physics parameters, like the average light neutrino mass, the right-handed week current admixtures and the heavy neutrino mass.

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