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Alpha Decay as a Strong Collective Phenomenon

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Abstract. A theory of the α-decay in terms of irreducible reaction amplitude of the α-particle formation in the four particle channel ($T_{4\to\alpha}$) is given. By inserting a new universal constant determined by the coupling interaction between the many body Fermi liquid state and the α-cluster state we have calculated the α-widths for some favoured and unfavoured α-transitions in the translead region. Good agreement with the experimental data is obtained.

1. Introduction

An analysis of the existent α-decay models[1-6] shows that the nucleus was treated therein as a Fermi gas object. This observation results from the way the α-particle has been treated in these models. Thus, the constituent nucleons of the α-particle either fly away independently[2-4] from the nucleus and subsequently gather together (cluster) to form an α particle, or the corresponding α-width formula uses in some way the first order perturbation approach.

The R-matrix theory[6] does not go along this line, but it encounters difficulties in defining the internal region nuclear states and the channel radius which are artificially introduced.

All these models lead to constantly lower values of the theoretical α-widths with respect to the experimental data[1,2,6].

The previous models propose themselves to transport to the continuum a group of four particles, from the mother nucleus, but they do not consider the mechanism of formation of the cluster. However, in our opinion it is more important to describe the α-clustering because, once we have done it, the α-decay process reduces to the well known problem of the barrier penetration.

The nuclear states involved in the α-decay models are described by the shell model with residual interactions where small momenta (of the order of $A^{-1/3}$) and energies participate. The same kind of interactions were also used in the first order perturbation theories[7-9] for the α-transition operator. However, one can hardly conceive, that such kind of interactions would be able to describe the clusterization where rather large momenta (of the order of $4^{-1/3}$) are involved. It is clear that some other kind of interactions must be responsible for the α-particle clustering

inside the nucleus and they cannot act in the frame of a Fermi gas model.

The nucleus is a system of strong interacting particles, where the mean distance between the nucleons is of the same order of magnitude as the range of the nuclear forces. This fact tells us that, the nucleus could be treated as a Fermi liquid with strong correlations among the nucleons.

The possibility of clusterization is equivalent to the coexistence of two different states in the nuclei. One state describes the nucleus as a Fermi liquid drop in which the quasi-particles are more or less uniformly distributed in its volume and another state describes two Fermi liquid fragments (α-particle and daughter nucleus) in the relatively weak interaction (a cluster state).

The decay of the cluster state is allowed or forbidded mainly by the barrier.

The presence of the above mentioned two states can be also found in the stable nuclei. Such a picture could explain why in the high energy nuclear reactions with different projectiles (p, d, α, light nuclei) one can observe a large yield of α-particle emission[2a-e].

It is natural to assume that the clusterization is a surface phenomenon. There are several arguments in favour of this assumption. At the nuclear surface the density is relatively low. The collisions of the α-particle with the neighbouring nucleons far inside the nucleus lead to a very improbable existence of the α-cluster in that region. The large incompressibility of the nuclear matter excludes the presence of innomogeneities like α-clusters far inside the nucleus. The α-scattering experiments are described by using a rather large imaginary part of the optical

potential which shows that the α-particles are melted inside the nucleus.

These arguments exclude the idea that the channel state wave functions have a resonant behaviour. However, the total wave function of the system must have a resonant character[1].

Assuming that the nucleus is a Fermi liquid, one can expect that the α-decay is a strong collective phenomenon taking place in two steps: the clusterization and the barrier penetration The clusterization process is a phase transition between the many body fermi liquid state to the α-cluster state. The corresponding transition operator has to force four quasinucleons distributed in the volume of the mother nucleus to occupy a small volume of the order of the α-particle volume. Such an operator must be determined by a new quantity, namely the irreducible amplitude of the α particle formation in the four particle channel[7]:

$T_{4\to\alpha}$ = 

(1)

An analogous situation occurs in the description of the pairing correlations in nuclei, where irreducible amplitudes[7] in the particle-particle or particle-hole channels are introduced.

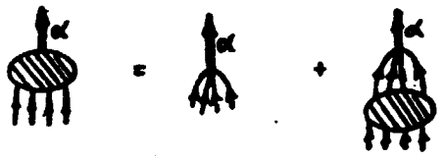
By assuming a model for the amplitude (1) we have calculated the α-widths for some favoured and unfavoured α-transitions between spherical nuclei in the translead region. An unique constant determined by fitting the favoured α-transition $^{210}Ra \to \alpha + {}^{206}Rn$ to the experiment was used.

2. Model for the Irreducible Amplitude of the α-particle Formation in the Four Particle Channel

Generally speaking, the amplitude $T_{4\to\alpha}$ can be extracted, in the framework of the many body theory, from the totality of diagrams of the four particle Green's function, excluding the terms, that involve the states around the Fermi sea. If $T^{(0)}_{4\to\alpha}$ is the zero-range irreducible amplitude

$$T^{(0)}_{4\to\alpha} = \quad (2)$$

we can write the equation for the $T_{4\to\alpha}$ as follows:



(3)

In the following we will propose a model for $T^{(0)}_{4\to\alpha}$ and neglect the second term in the eq. (3). The amplitude $T^{(0)}_{4\to\alpha}$ must have an universal character, determined by the properties of the nuclear matter (through the density dependence), i.e. by the states deep inside the Fermi sea:

$$T^{(0)}_{4\to\alpha} \simeq \varkappa(\vec{\ell}n\ \nabla)\rho(\vec{R})\delta(\vec{\xi}_1)\delta(\vec{\xi}_2)\delta(\vec{\xi}_3)\ell \quad (4)$$

Here $\vec{\ell}n$ is an unit normal to the nuclear surface, ∇ is the gradient operator,

$$\rho(\vec{R}) = \langle A+4|\hat{\rho}(\vec{R})|A+4\rangle;\ \hat{\rho}(\vec{R}) = \sum_{\lambda\lambda'} \varphi^*_\lambda(\vec{R})\varphi_{\lambda'}(\vec{R})a^+_\lambda a_{\lambda'} \quad (5)$$

in the nuclear matter density, where $|A+4\rangle$ stands for the initial state wave function, φ_λ - for the single particle wave function and $a^+_\lambda(a_\lambda)$ - for the fermion creation (absorption) operator of a nucleon in (from) the state $|\lambda\rangle$. $\vec{\xi}_i$ are the internal Jacobi coordinates of the α-particle. The $\delta(\vec{\xi}_i)$ functions describe the

packing process of the four nucleons in a small volume of the order of the α-particle volume. The $\mathcal{L}$ operator selects the terms containing two protons and two neutrons from among the four fermion orbitals.

The constant κ must have an unique value for all the α-transitions.

Our model is still in a preliminary stage, mainly because we have neglected the second term in the eq. (3). By solving the eq. (3) one could also attempt to calculate the α-decay energy, but this problem is as complicated as the similar problem of the binding energy. Therefore in our calculations the experimental α-decay energies are preferred.

The amplitude (4) in the second quantization form can be written as follows:

$$T^{(0)}_{A\to a} = \sum_{S_a S_A j\mu} T^{j}_{S_A S_a}(R)\, b^{+}_{S_a j\mu}\, A_{S_A j\mu} \tag{6}$$

Here

$$A_{S_A j\mu} = ((a_\nu a_{\nu'})_{j_p} (a_\omega a_{\omega'})_{j_n})_{j\mu} \tag{7}$$

where $\nu\nu'(\omega\omega')$ stand for the s.p. proton (neutron) orbitals (e.g. $\nu \equiv (n_\nu \ell_\nu j_\nu)$),

$$\langle \hat{R}, \xi_a | b^{+}_{S_a j\mu} | 0 \rangle = ((Y_{\ell_a}(\hat{R}) \psi_{S_\xi}(\vec{\xi}_1\vec{\xi}_2\vec{\xi}_3))_{L_a} \chi_{S_p S_n S_a}(\hat{S}_1\hat{S}_2\hat{S}_3\hat{S}_4))_{j\mu} \tag{8}$$

is the α-cluster wave function, where

$$\psi_{S_\xi m_\xi}(\vec{\xi}_1\vec{\xi}_2\vec{\xi}_3) = \langle b\vec{\xi}_1, b\vec{\xi}_2, b\vec{\xi}_3 | n_{\xi_1}\ell_{\xi_1} n_{\xi_2}\ell_{\xi_2}(\ell_{\xi_{12}}) n_{\xi_3}\ell_{\xi_3}; \ell_\xi m_\xi \rangle \tag{9}$$

is the spatial part and

$$\chi_{S_p S_n S_a}(\hat{S}_1\hat{S}_2\hat{S}_3\hat{S}_4) = \langle \hat{S}_1\hat{S}_2\hat{S}_3\hat{S}_4 | \tfrac{1}{2}\tfrac{1}{2}(S_p)\tfrac{1}{2}\tfrac{1}{2}(S_n), S_a \sigma_a \rangle \tag{10}$$

is the spin part.

The quantities $T^{J}_{S_a S_\beta}(R)$ are the reduced matrix elements $\langle S_a j \| T^{(0)}_{A \to a} \| S_\beta j \rangle$, when the eq. (4) is used.

3. Alpha Decay Width

Starting from the eq. (14) of ref.[1] (or the eq. (15) of ref.[8]) the complicated transition R-operator can be factorized by inserting the correlations in both the channel and initial states wave functions:

$$R = \Omega^{(A+a)\dagger} T_{A \to a} \Omega^{(A+A)} \tag{11}$$

The α-decay width can be written[1] as follows:

$$\Gamma_a = 2\pi \sum_c |\langle A+a, c | T_{A \to a} | A+A \rangle|^2 =$$

$$= 2\pi\hbar^2 \sum_c \Big| \sum_{c'} \Omega^{(A+a)}_{cc'} \int_0^\infty dR\, u_{c'}(R) \rho'(R) g^{S_i j_i}_{c'}(R) \Big|^2 \tag{12}$$

where the following notations have been used:

$$|S_j \mu\rangle = \phi_{S_j \mu} = \Omega^{\dagger}_{S_j \mu} |\tilde{0}\rangle \tag{13}$$

for the nuclear state wave functions,

$$|A+a, c\rangle \equiv n |u_c\rangle \equiv \sum_{c'} \Omega^{(A+a)}_{cc'} \mathcal{A} \Big(\frac{1}{R} u_{c'}(R) \big(b^{\dagger}_{S_a \ell_a} \Omega^{\dagger}_{S_\beta j_\beta} \big) \Big) |\tilde{0}\rangle \tag{14}$$

for the correlated channel wave function, in which c stands for the group $(S_a, \ell_a, j_\beta, S_\beta)$ of quantum numbers and $\mathcal{A}$ for the anti-symmetrization operator of the channel partners,

$$|A+A\rangle \equiv \Omega^{(A+A)} |\mathrm{USEC}\rangle \equiv \Omega^{\dagger}_{S_i j_i \mu_i} |\tilde{0}\rangle \tag{15}$$

for the initial state wave function,

$$g^{S_i j_i}_{c}(R) = R \langle \phi_a \big(Y_{\ell_a} \phi_{S_\beta j_\beta} \big)_{j_i \mu_i} | \delta(\vec{\xi}_1) \delta(\vec{\xi}_2) \delta(\vec{\xi}_3) | \phi_{S_i j_i \mu_i} \rangle =$$

$$= \sum_{S_A} R \langle S_a \ell_a \| \delta(\vec{\xi}_1) \delta(\vec{\xi}_2) \delta(\vec{\xi}_3) \| S_A \ell_a \rangle \times$$

$$\times \langle \vec{\theta} | [(a_{S_\beta j_\beta} A_{S_\alpha \ell_\alpha})_{j_i} \cdot a^+_{S_i j_i}] | \vec{\theta} \rangle \qquad (16)$$

We also used the eqs. (4) and (6) with

$$T^{\ell_\alpha}_{S_\alpha S_\alpha}(Q) = \rho'(R) \langle S_\alpha \ell_\alpha \| \delta(\vec{\xi}_1) \delta(\vec{\xi}_2) \delta(\vec{\xi}_3) \| S_\alpha \ell_\alpha \rangle , \qquad (17)$$

where $\rho'(R)$ is the first derivative of the nuclear density (5). The channel radial wave function $u_c(R) \equiv u_{cc}(R) \equiv u_c$ can be obtained from the total function (14) which is a solution of the following scattering eq.

$$(\epsilon - P\mathcal{H}P) P|A\alpha\rangle = 0, \qquad (18)$$

where

$$P = nN^2n^+, \quad P+Q = 1, \quad P^2 = P, \quad PQ = 0 = QP, \quad Q^2 = Q \qquad (19)$$

$$\mathcal{H} = H + HQ \frac{1}{\epsilon - H} QH \qquad (20)$$

and[10,18]

$$N^{-2} = 1 - K = n^+n \qquad (21)$$

From the eq. (18) we can extract the eq.

$$(\epsilon - Nn^+\mathcal{H}nN)|\tilde{u}_c\rangle = 0 \qquad (22)$$

where we have defined a new channel radial wave function

$$|\tilde{u}_c\rangle = N^{-1}|u_c\rangle \qquad (23)$$

with the same normalization condition

$$\langle \tilde{u}_c | \tilde{u}_{c'} \rangle = \delta(\epsilon - \epsilon') \qquad (24)$$

as for the channel functions (14).

It is easy to see that the renormalized Hamiltonian $Nn^+\mathcal{H}nN$ from the eq. (22) can be obtained by a folding procedure[9] including the exchange terms. We can obtain, afterwords the channel radial wave function $u_c(R)$, by using the integral eq. (23) and the N operator[10,18] from the eq. (21).

4. Numerical calculations

To compute the overlap integral $g_c^{S_i j_i}(R)$ (16) and the densities $\rho(R)$ (5), we need the structure of the initial and final nuclear states (13). For the favoured (g.s - g.s) α-transitions in the double even nuclei we have used the BCS-wave functions. It is known that the pairing correlations are very important in α-decay[11]. We included them by the following procedure[13-15]. The level scheme is recalculated for each nucleus, including all the bound levels and the quasibound levels with the widths up to 100 keV. Then, the pairing coupling constants G_p and G_n are fitted to reproduce the experimental pairing energies[12].

It must be noticed that excluding the quasibound levels the theoretical values decrease by a factor of ≈ 2[15].

As an example of unfavoured transition we have choosen the decay of ^{210}Bi to ^{206}Tl, because these nuclei differ from double magical ^{208}Pb by only one proton and one neutron and thus, their structure can be fairly well described. We used the nuclear structure from the ref.[16], where the configuration mixing is introduced to reproduce the low-lying level energies and the corresponding gamma and beta transitions.

To compute the radial channel wave function we neglected the exchange terms in the folding potential (i.e. $N \approx 1$[10,16]). Thus, for the potential occuring in the renormalized Hamiltonian of eq. (22) the direct term obtained by folding an Yukawa-type effective interaction has been analysed:

$$V_{Fold}(\vec{R}) = \int d^3x_1 d^3x_2 \rho_1(\vec{x}_1)\rho_2(\vec{x}_2)\sum_S \kappa_S(E)\,\frac{e^{-\mu_S|\vec{x}_1 - \vec{x}_2 - \vec{R}|}}{|\vec{x}_1 - \vec{x}_2 - \vec{R}|} \qquad (25)$$

$$\cdot \sum_S 16\pi^2 \varkappa_S(E) \int_0^\infty x_1^2 dx_1 \rho_1(x_1) \int_0^\infty x_2^2 dx_2 \rho_2(x_2) F^{(\mu)}(x_1 x_2 R)$$

where

$$F^{(\mu)}(x_1 x_2 x_3) = \{\operatorname{sign} Z_1 (1 - e^{-\mu|Z_1|}) + \operatorname{sign} Z_2 (1 - e^{-\mu|Z_2|}) +$$

$$+ \operatorname{sign} Z_3 (1 - e^{-\mu|Z_3|}) - (1 - e^{-\mu|Z|})\}(4\pi^2 x_1 x_2 x_3)^{-1}$$

with $Z_1 = -x_1 + x_2 + x_3$, $Z_2 = x_1 - x_2 + x_3$, $Z_3 = x_1 + x_2 - x_3$,

$$Z = x_1 + x_2 + x_3$$

$$\varkappa_0(E) = e^2;\ \varkappa_1(E) = 1528{,}75 \text{ Mev fm};\ \varkappa_2(E) = -784{,}4 \text{ MeV fm}$$

$$\mu_0 = 0;\ \mu_1 = 4 \text{ fm}^{-1};\ \mu_2 = 2.5 \text{ fm}^{-1}$$

The parameters $\varkappa_S(E), \mu_S$ were taken from ref.[10]. They were obtained by fitting the even-state G-matrix elements of the Reid interaction and assuming the odd-state interection of a purely OPEP type. No imaginary part has been used. The case $S = 0$ $(\varkappa_0, \mu_0)$ corresponds to the Coulomb part of the folding potential.

Unfortunately, the calculations of the channel radial function with the potential (26) are very cumbersome, therefore we compared the folding potential with the potential obtained by summing two single proton and two single neutron (fig. 1) shell model potentials:

$$V_4 = 2V_p + 2V_n \tag{26}$$

Assuming that the exchange terms from the eq. (24) lead to the decrease[38,40] of the difference between the folding potential and the potential (26) in the surface region, in the calculations we have used the potential $-V_0$.

In the figs. 2 and 3 we have plotted the overlap integral (16), the derivative (ρ') of the nuclear density of the initial

nucleus (5) and the radial channel wave function for the favoured ($^{210}Ra \rightarrow \alpha + {}^{206}Ra$) and unfavoured ($^{210}Bi$ (1-) $\rightarrow \alpha + {}^{206}Tl$ (1-))α-transitions using the set of Woods Saxon parameters taken from ref.[17]. These functions were also calculated with the s.p. Woods-Saxon's parameters obtained from the scattering experiments[20] and a small difference is obtained. This is shown in fig. 2 also. From these figures we conclude that the integrand from the eq.(12) is relatively large in the surface region only.

The calculated ratios $\frac{\Gamma_{exp}}{\Gamma_{theor}}$ for some favoured and unfavoured α-transitions are given in table 1. The value of the universal constant κ from the eq. (4) fitted to the favoured transition $^{210}Ra \rightarrow \alpha + {}^{206}Ra$ is equal to $1.044 \cdot 10^{7}$ Mev $fm^{1/2}$. We can see that the calculated α-widths for both favoured and unfavoured α-transitions are in agreement with the experimental data. Here we have choosen the α-transition, where the involved nuclear states have the best known structure. The calculations in the frame of the independent particle model for some Po and Ra isotopes lead to values of the ratios $\Gamma_{exp}/\Gamma_{theor}$ of the order of 100. For these isotopes nuclear correlations are hard to be introduced. The simple pairing model doesnot work. The energy gap practically vanishes. Pairing multipole correlations[21] are at least needed to describe the nuclear states of these isotopes. The above mentioned discrepancy has also been observed in the previous models[2]. Thus, we have eliminated from the beginning such cases from our statistics.

The remaining fluctuations (see table 1) are not large and they could be removed either by taking into account higher order terms in the eq. (3) and/or by including the exchange terms in the folding potential.

Of course, calculations in the actinide region have to be

done in addition. Such calculations are on the way.

5. Conclusions

In this work we have proposed a new picture to describe the α-decay phenomenon. First, we assume the coexistence of two kind of states in the nuclei: the fermi liquid state and the cluster state. The α-decay phenomenon takes place in two steps: the clusterization and the barrier penetration. The clusterization process is a (phase) transition between the Fermi liquid state and a α-cluster state. The barrier penetration process is practically the way the α-cluster state decays.

It is assumed that for practically all nuclei the Fermi liquid component in the structure of the nuclear states is dominant, therefore for the α-active nuclei, the clusterization process is the decisive step in the α-decay.

Our model is in a preliminary stage. The second term in the eq. (3) has to be also analysed, especially to study the α-decay energy. In order to be self-consistent we should also calculate the α-decay energy and obtain the α-decay width on this basis.

The exchange part of the folded potential might also give a significant contribution to the α-decay width.

Such calculations are on the way, but they lead to complicated expression for the potential, hard to be computed.

A coupled channel analysis (especially for deformed nuclei) has to be done, to obtain a realistic radial α-channel function.

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Table 1

Mother nucleus	A	$I_i^{\pi_i}$(E(Mev)) → $I_f^{\pi_f}$(E(Mev))	$\Gamma_{exp}/\Gamma_{theor}$
Ra	208	0^+(g.s) → 0^+(g.s)	0.56
	210	0^+(g.s) → 0^+(g.s)	1.00
	212	0^+(g.s) → 0^+(g.s)	1.56
	214	0^+(g.s) → 0^+(g.s)	2.95
Rn	202	0^+(g.s) → 0^+(g.s)	0.31
	206	0^+(g s) → 0^+(g.s)	1.60
	210	0^+(g.s) → 0^+(g.s)	1.56
Bi	210	1^-(g.s) → 1^-(0.304)	1.54
		1^-(g.s) → 2^-(0.266)	1.26
		9^-(0.205)→1^-(0.304)	7.10
		9^-(0.265)→2^-(0.266)	0.2

Caption to the table 1

The calculated $\Gamma_{exp}/\Gamma_{theor}$ ratios with $\varkappa = 1.044 \cdot 10^7$ Mev $fm^{1/4}$ and Woods-Saxon parameters from ref.[20].

Caption to the figures

Fig. 1. The direct term of the folded potential and V_c obtained with the set of Woods-Saxon parameters taken from refs. refs.[17,20] for the $\alpha + {}^{208}Pb$ channel.

Fig. 2. $g_c^{S_i,j_i}$, ρ' and U_c - functions obtained with the set of Woods-Saxon parameters taken from refs.[17,20], for the α-decay of ^{210}Ra.

Fig. 3. The same as in the fig. 2, but with the Woods-Saxon parameters taken from ref. for the α-decay of ^{210}Bi (1^-) to ^{206}Tl (1^-).

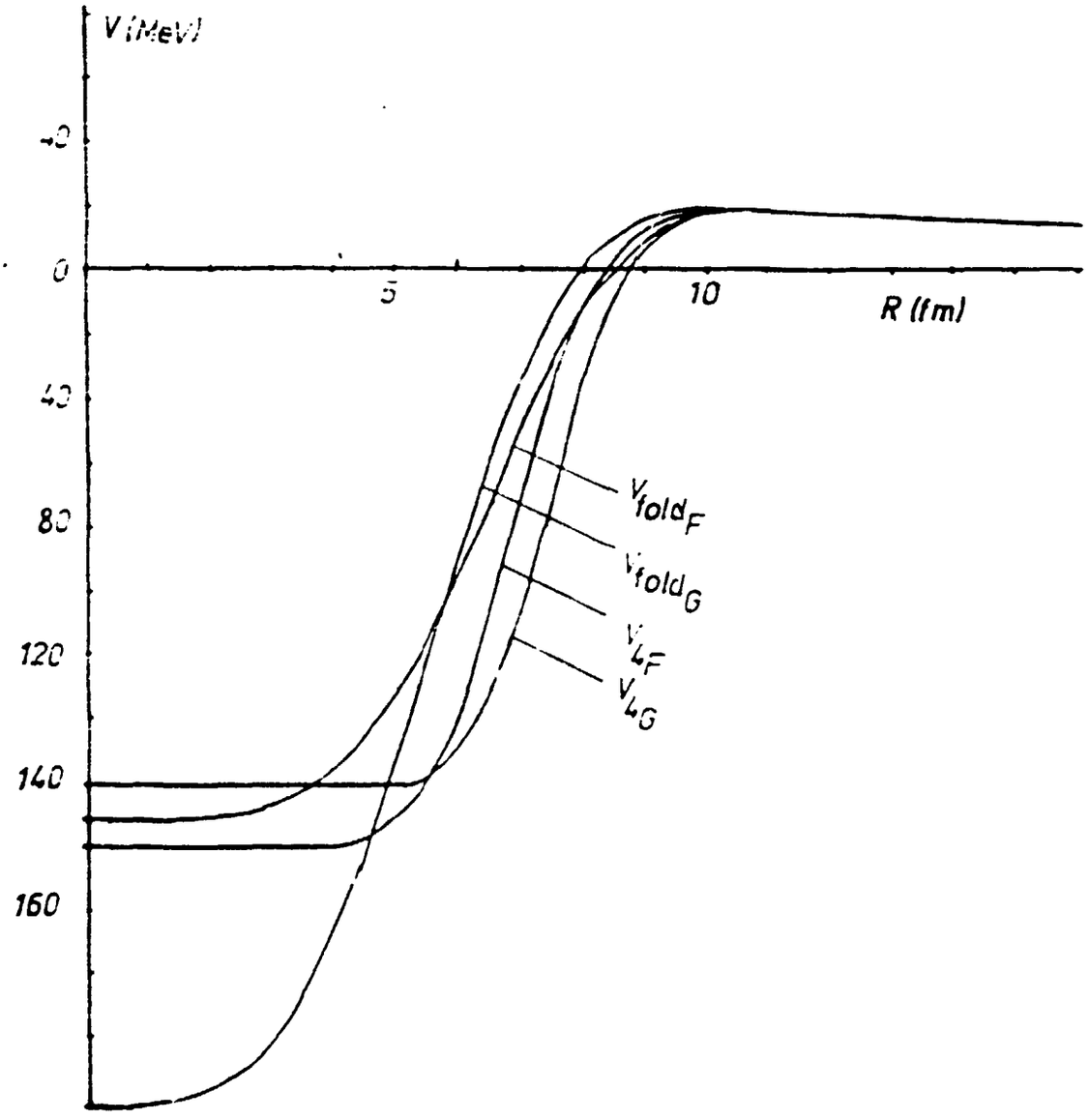

Fig. 4

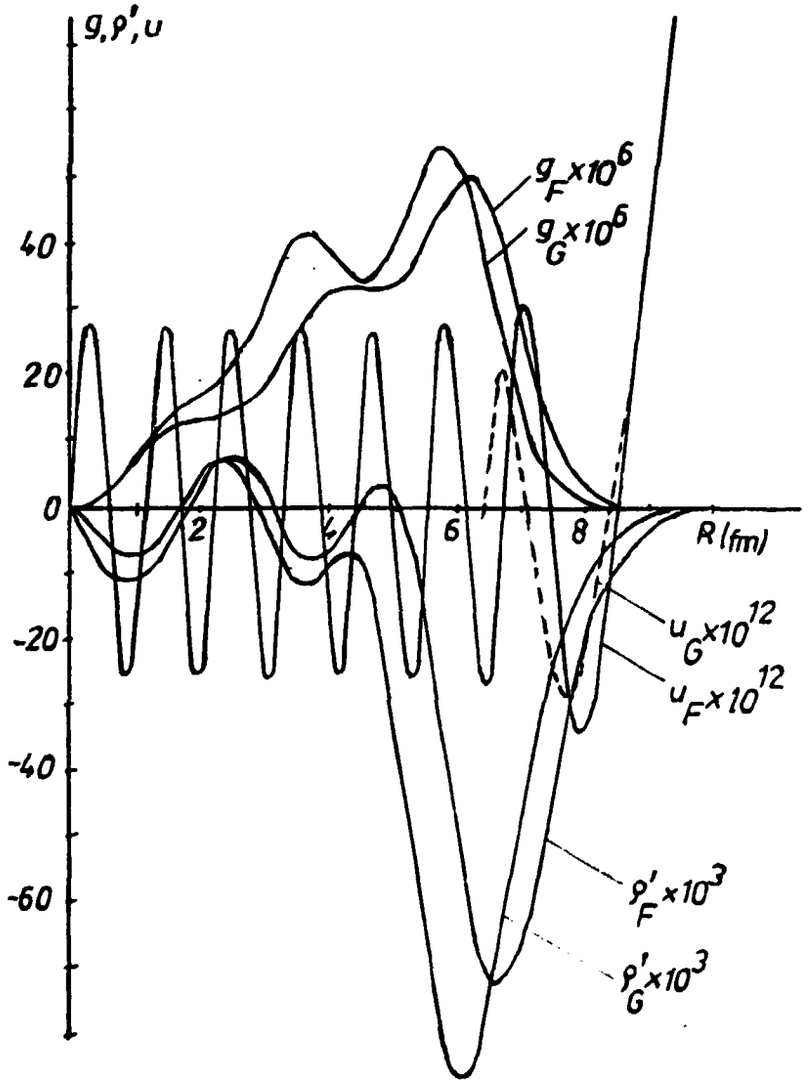
g, ρ', u
40
20
0
-20
-40
-60
2
4
6
8
R (fm)
$g_F \times 10^6$
$g_G \times 10^6$
$u_G \times 10^{12}$
$u_F \times 10^{12}$
$\rho'_F \times 10^3$
$\rho'_G \times 10^3$

Fig 2

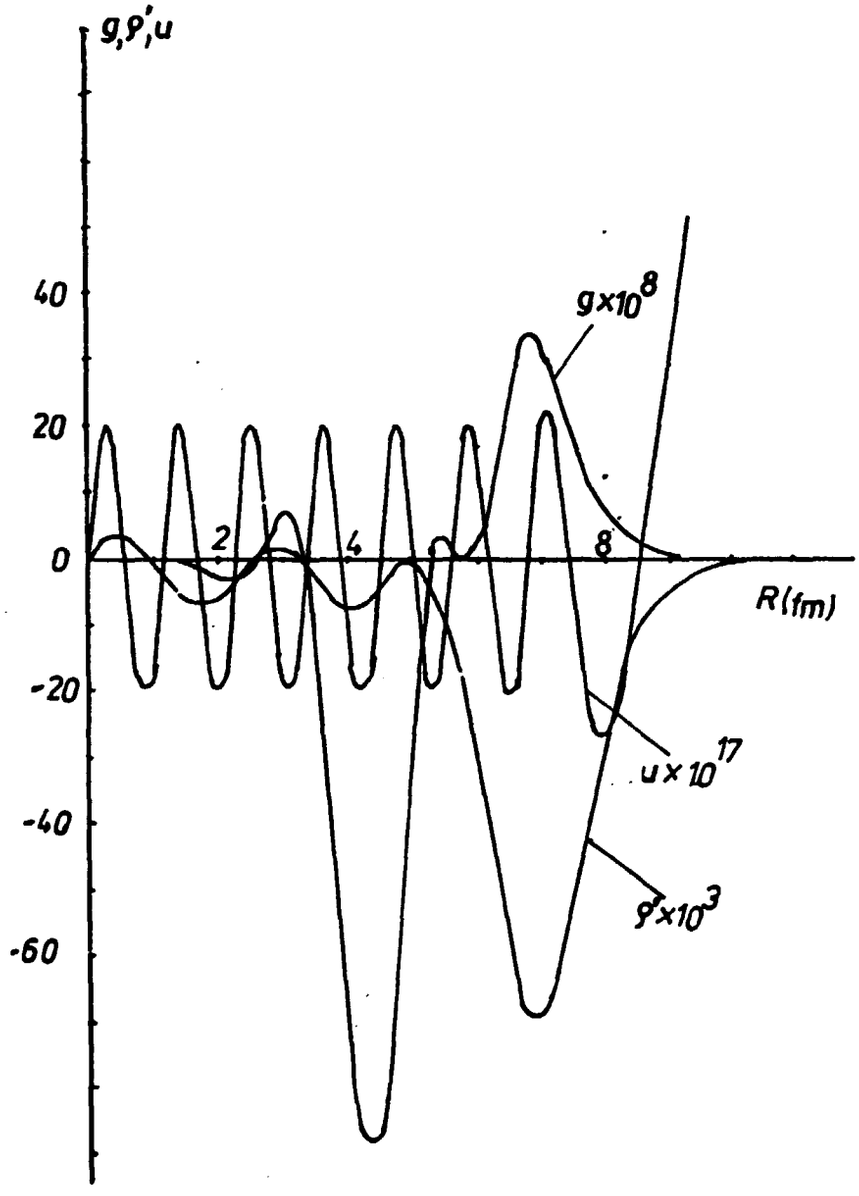

Fig. 3

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