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ON THE EXISTENCE OF COHERENT ROTATIONAL STATES IN NUCLEI

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INTRODUCTION

If two heavy ions scatter with an energy of the relative motion below, but close, the Coulomb barrier height, the main interaction between them is the Coulomb interaction, which gives rise to a hyperbolic trajectory of the relative motion. Due to the finite extension of the interacting nuclei and the presence of the nuclear interaction at short distances, the electric and nuclear fields may cause excitations in both projectile and target heavy ions. These excitations are of the collective - rotational and vibrational character. For nuclei which possess an equilibrium shape which is not spherically symmetric, the lowest excited states form a rotational band, generated by rotating the nucleus to states of higher angular momentum. In the simplest picture the intrinsic structure of the deformed nucleus is independent of the rotation and the spectrum will be that of a rigid rotor with constant moment of inertia J $(E_{rot} = \hbar\omega I(I+1)$, where I is the angular momentum). Among the excitations there can be found the low-lying vibrational states classified into quadrupole, octupole, etc., modes, giving rise to a harmonic spectrum with equidistant level spacing determined by the basic frequency. Coulomb excitations below the Coulomb barrier are specifically suited to excite the rotational and low-frequency vibrational modes. Because of slow variation of the Coulomb plus nuclear field in space, it can only excite modes of low multiplicity.

For bombarding energies below, but close to, the Coulomb barrier, excitation of a rotational state enters through the wave function

$$|\psi(t)\rangle = \sum_{IKM,\nu} a_{IKM,\nu}\, e^{-iE_{IK\nu} t/\hbar}\, |IKM,\nu\rangle , \qquad (1)$$

which depends on the excitation amplitudes $a_{IKM,\nu}$ in a coherent fashion. The eigen-states $|IKM,\nu\rangle$ are defined in an adiabatic model as follows:

$$|IKM,\nu\rangle = D^{I}_{MK}(\theta_i)\chi_{K\nu} , \qquad (2)$$

where $D^{I}_{MK}(\theta_i)$ are the normalized to unity eigen-states of the rigid rotor and $\chi_{K\nu}$ describes the intrinsic motion.

The wave function (1) is a coherent rotational states (CRS) if the amplitudes $a_{IKM,\nu}$ have equidistant phases and their absolute values are peaked around a mean value of the angular momentum (I) of rotation [1]. Coherent rotational states were studied to see the correspondence between classical and quantum mechanical rotation.

The properties of CRS are similar formally to those of coherent states in quantum optics [2]. The simplest example is the system of photons, all of which have the same momentum and polarization. The coherent state in this case is nothing else but the state to be identified with the coherent plane electromagnetic wave in classical electromagnetism:

$$|\alpha\rangle = \sum_{n=0}^{\infty} \exp(in\alpha)\, w_n^{1/2}\, |n\rangle \,, \tag{3}$$

where

$$|n\rangle = (n!)^{-1/2} (b^+)^n |0\rangle \tag{4}$$

is the normalized eigen-function belonging to the eigen-value n of the photon number operator b^+b (b^+, b are the creation and destruction operators of photons). The coherent state (3) is the wave packet obtained by superposing the $|n\rangle$ with the same phase factor (α) and w_n is a probability distribution, which has a sharp peak at a macroscopic photon number N with a width of the order of $N^{1/2}$. The coherent state (3) is used in quantum optics to represent the state of laser light. It can be applied to other Bose systems as, e.g., for rigid rotors [1]. Assuming that in the heavy ion collision in which one heavy ion is deformed the Coulomb interaction generates wave packets of rotational states (WPRS) of the form (1) in which the amplitudes show the coherence properties [3-5].

Here and in the following we shall call such WPRS CRS.

The decay of the CRS has been studied in refs. [5-7], where it was shown that when the target nucleus has an ideal rotational band (i.e., $E_{rot} = \hbar\omega I(I+1)$), the coherence of the state recovers periodically and this gives rise to strong pulses in the spontaneous gamma and alpha decay probabilities. When the rotational band is not ideal, then the decay probability is quasiperiodic in time.

In the preceding paper [5] we have shown that the inclusion of the nuclear part of the interionic interaction potential (within the folding procedure) does not destroy the coherence properties of the CRS-amplitudes, even if the absorbtive part is few per cent of the real one.

The purpose of this paper is to replace the absorbtive part of the interionic interaction potential by coupling to some

internal degrees of freedom (e.g., to the β-vibrational phonon) and to study the influence of this coupling on the CRS-amplitudes. The β-vibrational band is the closest rotational band to the ground state band and it is easier to be excited in the collision process as compared to the other (gamma, octupole,...) bands.

We show that the coherence of CRS-amplitudes is not destroyed if coupling to β-band, and if some conditions are satisfied.

Thus the conclusion is that the intermediate excitations are weak and still preserve the coherence properties of CRS-amplitudes when including nuclear interaction among the heavy ions, conclusion that leads to the idea that the CRS may be produced in heavy ion collision.

II. COUPLING TO β-VIBRATIONAL BAND INTERACTION POTENTIAL

Here we assume as before [6] the sudden approximation and the folding potential [8] acting between the two heavy ions in the collision process:

$$V(\vec{R}) = V^{coul}(\vec{R}) + V^{nucl.}(\vec{R}) =$$

$$= \sum_{s\lambda_1\lambda_2\lambda_3} \sum_m \kappa_s (4\pi)^{3/2} \hat{\lambda}_1\hat{\lambda}_2\hat{\lambda}_3 \begin{pmatrix} \lambda_1 & \lambda_2 & \lambda_3 \\ 0 & m & -m \end{pmatrix} \begin{pmatrix} \lambda_1 & \lambda_2 & \lambda_3 \\ 0 & 0 & 0 \end{pmatrix} \times \tag{5}$$

$$\times Y_{\lambda_3 m}(\hat{R})\, D^{\lambda_2}_{m0}(\hat{e}) \int_0^\infty x_1^2 dx_1\, \rho_{\lambda_1}(x_1) \int_0^\infty x_2^2 dx_2\, \hat{\rho}_{\lambda_2}(x_2)\, F^{(\mu_s)}_{\lambda_1\lambda_2\lambda_3}(x_1 x_2 R)$$

in which $F^{(\mu)}_{\lambda_1\lambda_2\lambda_3}(x_1 x_2 R)$ is approximated with:

$$F^{(\mu)}_{\lambda_1\lambda_2\lambda_3}(x_1 x_2 R) \cong \Theta(R - x_1 - x_2)\, i\mu\, h^{(+)}_{\lambda_3}(i\mu R)\, j_{\lambda_2}(i\mu x_2)\, j_{\lambda_1}(i\mu x_1) .$$

Here the notation is the same as in refs. [5,8] except

$$\hat{\rho}(\vec{x}_2) = \sum_{\lambda_2} \hat{\rho}_{\lambda_2}(x_2) Y_{\lambda_2 0}(\hat{x}_2) = \sum_{\nu\tau\nu'\tau'} \phi^*_{\nu\tau}(\vec{x})\phi_{\nu'\tau'}(\vec{x})\, a^+_{\nu\tau} a_{\nu'\tau'} =$$

$$= \sum_{\nu\nu'} \phi^*_{\nu+}(\vec{x})\phi_{\nu'+}(\vec{x}) \{ v_{\nu\nu'} B(\nu\nu') + \frac{1}{\sqrt{2}} u_{\nu\nu'} (A^+(\nu\nu') + A(\nu\nu')) \} + \tag{6}$$

$$+ \sum_{\nu\nu'} \phi^*_{\nu+}(\vec{x})\phi_{\nu'-}(\vec{x}) \{ v_{\nu\nu'} \bar{B}(\nu\nu') + \frac{1}{\sqrt{2}} u_{\nu\nu'} (\bar{A}^+(\nu\nu') + \bar{A}(\nu\nu')) \} ,$$

where $\phi_{\nu\tau}(\vec{x})$ are the s.p. w.f. generated by a static deformed potential [9,10]. The quantities $u_{\nu\nu'}$ and $v_{\nu\nu'}$ are of the form

$$u_{\nu\nu'} = u_\nu v_{\nu'} + v_\nu u_{\nu'} , \tag{7}$$

$$v_{\nu\nu'} = u_\nu u_{\nu'} - v_\nu v_{\nu'} \tag{8}$$

in which u_ν, v_ν are the Bogolubov-Valatin transformation coefficients [10]:

$$a^+_{\nu\tau} = u_s \alpha^+_{s-\sigma} + \sigma v_s \alpha_{s\sigma} , \tag{9}$$

and

$$B(ss') = \sum_\sigma \alpha^+_{s\sigma} \alpha_{s'\sigma} , \tag{10}$$

$$\bar{B}(ss') = \sum_\sigma \sigma \alpha^+_{s-\sigma} \alpha_{s'\sigma} , \tag{11}$$

$$A^+(ss') = \frac{1}{\sqrt{2}} \sum_\sigma \sigma \alpha^+_{s'-\sigma} \alpha_{s\sigma} , \tag{12}$$

Keeping the same conditions as in ref. [5] for the CRS-amplitudes (i.e., the sudden approximation and the quadrupole part of the coupling interaction (5)) we extract the part of the potential that excites the β-band:

$$V_\beta(\vec{R}) \cong \sum \kappa_s (4\pi)^{3/2} Y_{20}(\hat{R}) D^2_{00}(\hat{\theta}) \int_0^\infty x_1^2 dx_1 \rho_0(x) \times$$
$$\times \int_0^\infty x_2^2 dx_2 \hat{\rho}_{2,\beta}(x_2) i\mu_s h_2^{(+)}(i\mu_s R) j_0(i\mu_s x_1) j_2(i\mu_s x_2) . \tag{13}$$

Here we have assumed that the projectile remains unexcited during the collision and the excitation of the target nucleus is performed due to $\hat{\rho}_{2,\beta}$-operator. We have assumed also that $\Theta(R-x_1-x_2)=1$ which is almost a good approximation, because the distance of the closest approach (R_{min}) is always greater than the sum of the two radii of the colliding partners.

III. STRUCTURE OF COHERENT ROTATIONAL STATES

According to the semiclassical description [3,4] of the formation of CRS, the amplitudes of CRS (see eq. 1) in the case of ground and β- bands ($K = 0$, $\nu = \beta$ or g) are given as follows:

$$a_{IMK,\nu} = \delta_{K0} < IM0,\nu \,|\, \hat{a}_1 \hat{a}_2 \,|\, 000, g > , \tag{14}$$

where

$$\hat{a}_1 = \exp\{-\frac{i}{\hbar} \int_{-\infty}^{+\infty} dt (V^{coul.}(t) + V^{nucl.}(t)) , \tag{15}$$

$$\hat{a}_2 = \exp\{-\frac{i}{\hbar} \int_{-\infty}^{+\infty} dt V_\beta(t) \} . \tag{16}$$

The expression for $< IM0, g | \hat{a}_1 | 000, g >$ is given in ref. [5]:

$$a_{IM,g} = < IM0, g | \hat{a}_1 \hat{a}_2 | 000, g > = < IM0, g | \hat{a}_1 | 000, g > =$$
$$= \delta_{M0} \frac{\hat{I}}{2} \int_{-1}^{+1} dx P_I(x) e^{-iA(x)} , \tag{17}$$

where

$$A(x) = \frac{Z_1 e^2 Q_0}{8\pi v a^2}(3x^2 - 1) + (v_0 - iw_0)(A_g + B_g(3x^2 - 1)) , \tag{18}$$

$$a = \frac{b}{2} = \frac{Z_1 Z_2 e^2}{2E} , \tag{19}$$

$$A_g = \sum_{s=1}^{2} \kappa_s \tilde{A}_{1s} \tilde{A}_{2s} \frac{2a\mu_s}{\hbar v} I_0(a\mu_s) , \tag{20}$$

$$B_g = \sum_{s=1}^{2} \kappa_s 2\sqrt{5\pi}\, \tilde{A}_{1s} \tilde{\mathfrak{M}}^{(2)}_{20,s}(g) \frac{2a\mu_s}{\hbar v} I_2(a\mu_s) , \tag{21}$$

$$\tilde{A}_{js} = \sqrt{4\pi} \int_0^\infty x^2 dx\, \rho_0^{(j)}(x) i_0(\mu_s x) , \tag{22}$$

$$\tilde{\mathfrak{M}}^{(j)}_{20,s}(g) = \int_0^\infty x^2 dx\, \rho_2^{(j)}(x) i_2(\mu_s x) = 2 \sum v_\nu^2 < \nu + | i_2 Y_{20} | \nu + > , \tag{23}$$

$$I_0(x) = K_0^{(+)}(x) e^{-x} , \tag{24}$$

$$I_2(x) = (-\frac{1}{2} K_0^{(+)}(x) + K_1^{(+)}(x) + \frac{1}{2} K_2^{(+)}(x)) e^{-x} , \tag{25}$$

$$i_\ell(x) = \sqrt{\frac{\pi}{2x}} I_{\ell+1/2}(x); \quad K_\ell^{(+)}(x) = \sqrt{\frac{2}{\pi x}} K_{\ell+1/2}(x) . \tag{26}$$

Here Q_0 is the quadrupole moment of the target; b, the distance of the closest approach; Z_i, the atomic numbers; V, the speed of the projectile; $\rho_\lambda(x)$, the multipoles of the nuclear density; κ_s, μ_s, V_0, W_0, the folding potential parameters [8]. I_ℓ and K_ℓ are the modified spherical Bessel and Hankel functions, respectively, $\langle \nu\tau | i_2 Y_{20} | \nu'\tau' \rangle$ is the s.p. matrix element of $i_2 Y_{20}$.

To calculate the $\langle \beta | \hat{a}_2 | g \rangle$ matrix element we use the RPA technique for the structure of the β-vibrational state:

$$|\beta\rangle = Q^+_\beta |g\rangle \quad (27)$$

with

$$Q^+_\beta = \frac{1}{2} \sum_{\nu\nu'} (\psi^\beta_{\nu\nu'} A^+(\nu\nu') - \phi^\beta_{\nu\nu'} A(\nu\nu')) \quad (28)$$

(see eq. 12) and:

$$\langle \beta | \hat{a}_2 | g \rangle = \langle g | Q_\beta e^{-iT} | g \rangle, \quad (29)$$

where

$$T = \frac{1}{\hbar} \int_{-\infty}^{+\infty} dt\, V_\beta(t). \quad (30)$$

Using the equation:

$$\langle g | [Q_\beta, T] | g \rangle \cong B_\beta (3x^2 - 1) \quad (31)$$

with

$$B_\beta = \sum_{s=1}^{2} \kappa_s 2\sqrt{5\pi}\, \tilde{A}_{1s} \tilde{\mathfrak{M}}^{(2)}_{20,s}(\beta) \frac{2a\mu_s}{\hbar v} I_2(a\mu_s), \quad (32)$$

in which

$$\tilde{\mathfrak{M}}^{(j)}_{20,s}(\beta) = \sum_{\nu\nu'} u_{\nu\nu'} (\psi^\beta_{\nu\nu'} + \phi^\beta_{\nu\nu'}) \langle \nu + | i_2 Y_{20} | \nu' + \rangle \quad (33)$$

we find within RPA that

$$\langle g | Q_\beta e^{-iT} | g \rangle \cong -i B_\beta (3x^2 - 1) e^{-B^2_\beta (3x^2-1)} \quad (34)$$

and finally

$$a_{IM0,\beta} = \langle IM0, \beta | \hat{a}_1 \hat{a}_2 | 000, \beta \rangle = -i\delta_{M0} \frac{\hat{I}}{2} B_\beta \int_{-1}^{+1} dx\, P_I(x)(3x^2-1) e^{-iA(x) - B^2_\beta (3x^2-1)^2}. \quad (35)$$

IV. ALPHA DECAY OF COHERENT ROTATIONAL STATES

The time derivative of the α-decay probability as a function of time has the following expression now (see also refs. [5,7]):

$$\frac{dw_{if}}{dt} = \frac{1}{2} \sum_{\substack{I_i I_i' \\ \nu_i \nu_i'}} \sum_{\substack{I_f I_f' \\ \nu_f \nu_f'}} \sum_L C^{L I_f I_i}_{0\,0\,0} C^{L I_f' I_i'}_{0\,0\,0} \times$$

$$\times |a_{I_f \nu_f} a_{I_i \nu_i} a_{I_f' \nu_f'} a_{I_i' \nu_i'}| \cos[(\Delta_{i'f'} - \Delta_{if})t + \Delta\phi_{i'f'} - \Delta\phi_{if}] \times$$

$$\times \{\Theta(\Delta_{if} + Q_\alpha) \Gamma^{1/2}_{\Delta_{if}+Q_\alpha}(I_i\nu_i, I_f\nu_f)_L \Gamma^{1/2}_{\Delta_{if}+Q_\alpha}(I_i'\nu_i', I_f'\nu_f')_L +$$

$$+ \Theta(\Delta_{i'f'} + Q_\alpha) \Gamma^{1/2}_{\Delta_{i'f'}+Q_\alpha}(I_i\nu_i, I_f\nu_f)_L \Gamma^{1/2}_{\Delta_{i'f'}+Q_\alpha}(I_i'\nu_i', I_f'\nu_f')_L\},$$

where

$$a_{I,\nu} = |a_{I,\nu}| e^{i\phi_{I\nu}},$$

$$\Delta_{if} \doteq \hbar\omega_{\nu_i} I_i(I_i+1) - \hbar\omega_{\nu_f} I_f(I_f+1),$$

$$\Delta\phi_{if} = \phi_{I_i\nu_i} - \phi_{I_f\nu_f},$$

$$Q_\alpha = E_{g.s.i.} - E_{g.s.f.} - E_{g.s.\alpha},$$

$$\Gamma^{1/2}_\epsilon (I_i\nu_i, I_f\nu_f)_L = (-)^L C^{I_i L I_f} F^{1/2}_L b_L(\nu_i\nu_f) \Gamma^{1/2}_\epsilon(g.s.; g.s.),$$

$$F_L = \frac{P_L(\epsilon, R_0)}{P_0(\epsilon, R_0)} \cong \exp\{-2.027 L(L+1) Z^{-1/2} A^{-1/6}\},$$

$$\Gamma_\epsilon(g.s.; g.s.) = 6.5819 \ln 2 \exp(\ln 10 (M - \frac{N}{\sqrt{\epsilon}})),$$

in which for ^{238}U: M = 53.2139, N = 145.95.

The relative wave amplitudes [11,12] are defined as follows:

$$b_L(\nu_i, \nu_f) = \frac{\sum_{L'} K_{LL'} M^{L'}_{\nu_i\nu_f}}{\sum_{L'} K_{0L'} M^{L'}_{gg}}$$

in which [13]

$$M^L_{\nu_i\nu_f} = \langle g | [\Omega(\nu_i), \Phi^+_{\nu_f L}] | g \rangle,$$

where

$$[\Omega(\nu_i),\Phi^+_{\nu_i L}] = \begin{cases} A^+_4 & \text{for } \nu_i = g,\ \nu_f = g, \\ [Q_\beta, A^+_4] & \text{for } \nu_i = \beta,\ \nu_f = g, \\ A_4 Q^+_\beta & \text{for } \nu_i = g,\ \nu_f = \beta, \\ [Q_\beta, A^+_4 Q^+_\beta] & \text{for } \nu_i = \beta,\ \nu_f = \beta, \end{cases}$$

and

$$A^+_4 = \frac{1}{4} \sum_{\substack{\nu\nu' \\ \pi\pi'}} \sum_{\substack{\omega\omega' \\ \sigma\sigma'}} A^{L0}_{\pi\pi'\sigma\sigma'} (\nu\nu'|\omega\omega')\, a^+_{\nu\pi} a^+_{\nu'\pi'} a^+_{\omega\sigma} a^+_{\omega'\sigma'} \quad .$$

Thus

$$M^L_{gg} = \sum_{\nu\omega} A^{L0}_{+-+-} (\nu\nu|\omega\omega)\, \xi_\nu \xi_\omega ,$$

$$M^L_{g\beta} = \sum_{\nu\nu'\omega} A^{L0}_{+-+-} (\nu\nu'|\omega\omega)\, f_{\nu\nu'} \xi_\omega ,$$

$$+ \sum_{\nu\omega\omega'} A^{L0}_{+-+-} (\nu\nu|\omega\omega')\, \xi_\nu f_{\omega\omega'} ,$$

$$M^L_{\beta g} = \sum_{\nu\nu'\omega} A^{L0}_{+-+-} (\nu\nu'|\omega\omega)\, g_{\nu\nu'} \xi_\omega ,$$

$$+ \sum_{\nu\omega\omega'} A^{L0}_{+-+-} (\nu\nu|\omega\omega')\, \xi_\nu g_{\omega\omega'} ,$$

$$M^L_{\beta\beta} = \sum_{\nu\nu'\omega\omega'} A^{L0}_{+-+-} (\nu\nu'|\omega\omega')\, g_{\nu\nu'} f_{\omega\omega'} ,$$

where

$$\xi_\nu = u_\nu v_\nu ,$$

$$f_{\nu\nu'} = v_\nu v_{\nu'} \psi^\beta_{\nu\nu'} - u_\nu u_{\nu'} \phi^\beta_{\nu\nu'} ,$$

$$g_{\nu\nu'} = v_\nu v_{\nu'} \phi^\beta_{\nu\nu'} - u_\nu u_{\nu'} \psi^\beta_{\nu\nu'} ,$$

and $K_{LL'}$ are the Fröman matrices.

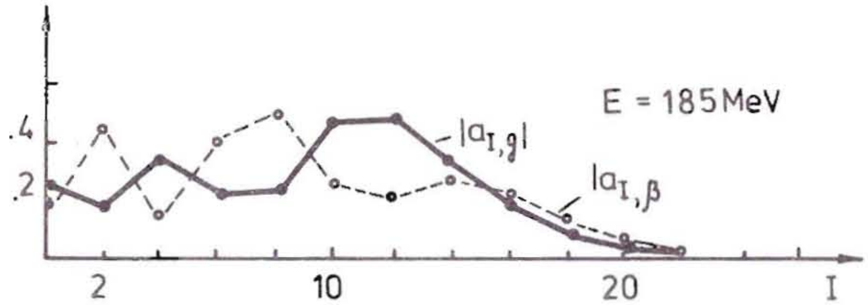

Fig. 1. The absolute values of the CRS amplitudes a_{Ig} and $a_{I\beta}$ for E = 185 MeV, taking into account the Coulomb and the real nuclear potential.

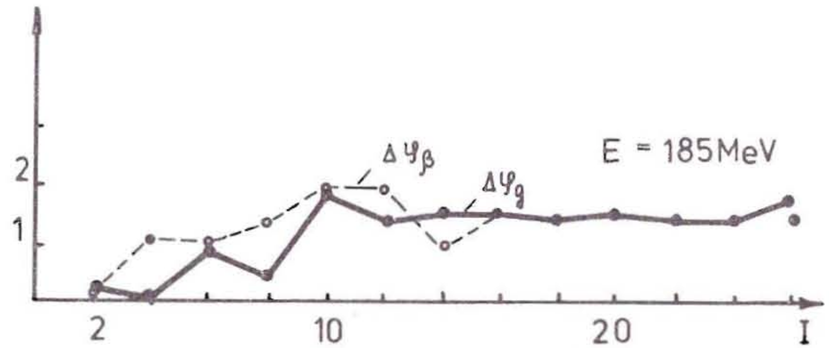

Fig. 2. The phase differences $\Delta\phi_{I\nu} = |\phi_{I,\nu} - \phi_{I-2,\nu}|$ for E = 185 MeV, taking into account the Coulomb and the real nuclear potential.

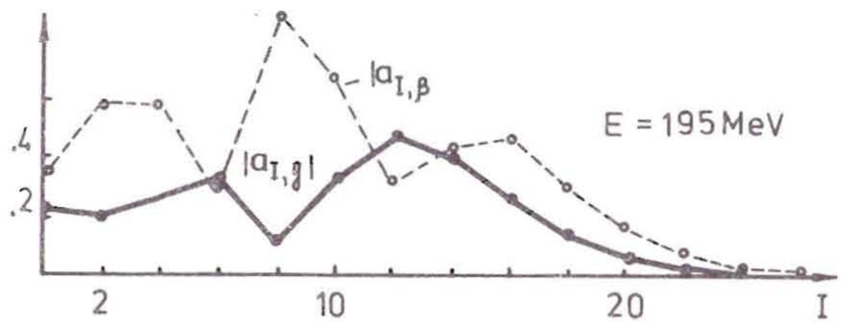

Fig. 3. The same as in fig. 1 but for E = 195 MeV.

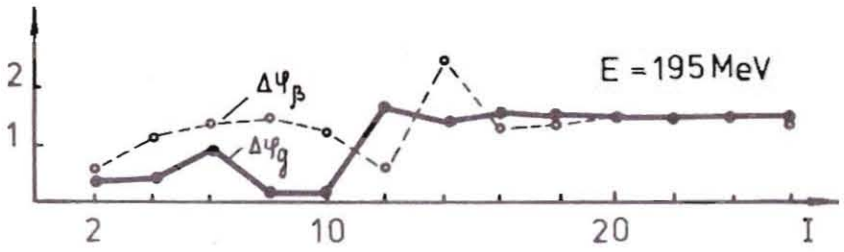

Fig. 4. The same as in fig. 2 but for E = 195 MeV.

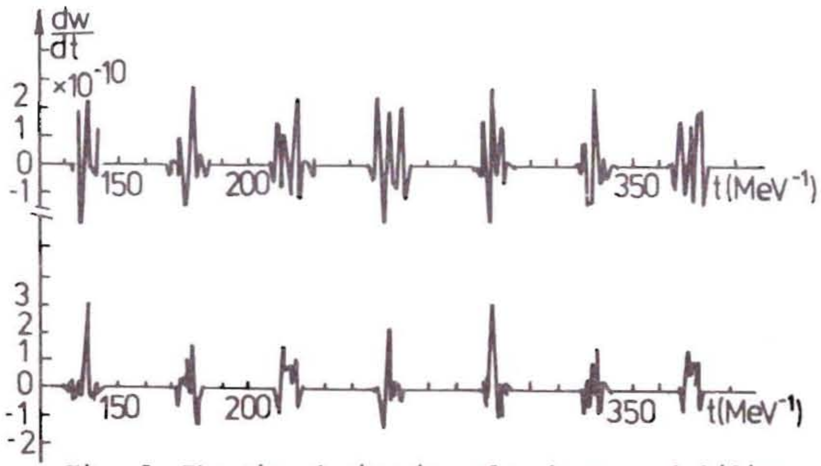

Fig. 5. The time derivative of α-decay probability for the cases when the stucture of CRS amplitudes is obtained with: a) Coulomb only; b) Coulomb plus nuclear potentials involving the β-vibrational band also.

V. NUMERICAL EXAMPLE

We present here the same numerical example as in refs./5-7/ and within the same conditions (see table 1 of ref./5/) which are well enough for the non-relativistic approximation and classical trajectories to be acceptable. The energy of the projectile E = 185 MeV (or 195 MeV) is 88% (or 92%) of the height of the Coulomb barrier.

In figs. 1-4 we have calculated the CRS-amplitudes determined by the Coulomb plus nuclaer folding interionic potential and with the coupling to β-vibrational band of ^{238}U-nucleus.

We see that the coherent properties of the CRS-amplitudes are preserved for both a_{Ig} and $a_{I\beta}$.

Fig. 5 shows the time derivative of the α-decay probability from CRS, which contain the CRS-amplitudes calculated in figs. 1,2. The periodicity of the pulses is the same as for the Coulomb case.

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Брындуш И., Карстою В., Думитреску О. E4-83-16

О существовании когерентных ротационных состояний в атомных ядрах

Изучается влияние бета-вибрационной полосы на структуру когерентных ротационных состояний, появляющихся при столкновении тяжелых ионов. Показывается, что когерентные свойства не исчезают, если учитывается бета-вибрационная полоса, когда энергия снаряда не превышает 75% высоты барьера и, поглощая часть потенциала, составляет несколько процентов от действительной части межионного потенциала.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

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On the Existence of Coherent Rotational States in Nuclei

The contribution of the β-vibrational band coupling in the structure of the coherent rotational states (CRS) formed in the heavy ion collision is analysed. It is shown that the coherence does not disappear when including virtual excitations to the members of the β-band if the energy of the projectile is not too high ($\approx$75% of the barrier's height) and if the absorbtive part is few per cent of the real part of the interionic potential.

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

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