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Nguyen Dinh Dang*

INFLUENCE OF THE SUPERFLUID PAIRING INTERACTION ON THE COLLECTIVE STATES IN THE FINITE TEMPERATURE RANDOM PHASE APPROXIMATION

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* Moscow State University

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The recent discovery of giant resonances built on states above the yrast line of highly excited nuclei in heavy ion fusion [1-3] and deep inelastic reactions [4] has sparked off a great interest in studying the properties of collective states in systems with large intrinsic energy and high spin. For these systems the microcanonical description employed to the states of the equilibrated nucleus with a given value of the temperature [5] can be applied. Thus, several approaches have been advanced in the framework of the finite temperature random phase approximation (FT-RPA) to investigate the temperature dependence of such nuclear characteristics as the location of giant multipole resonances (GMR), their strength distributions, the energy weighted sum rule (EWSR) [6-13]. The FT-RPA equations in these studies have been derived in the Hartree-Fock or Hartree-Fock-Bogolubov theory and the case of superfluid systems has been considered by solving the temperature dependent gap equations. It is well known that in superfluid systems there is a collapse of the pairing gap, which is complete for a critical temperature due to the phase transition from the superfluid state to the normal one* [6,12,14]. However, in the investigations of the temperature dependence of giant resonances the temperature region near this critical point has not yet been considered. It is the aim of the present letter to study the behaviour of giant resonance characteristics at finite temperature with emphasis on the critical temperature region for the superfluid system and to investigate how the change of the superfluid gap in dependence on temperature can influence the location of the giant resonances as well as their strength distributions.

We investigate the temperature dependence of dipole and quadrupole modes in ^{58}Ni-nucleus for which the superfluid pairing interaction is large at zero temperature. For comparison with the case of closed-shell nuclei we show also as an example the calculations of the isoscalar quadrupole resonance in ^{208}Pb-nucleus. We are going to employ the quasiparticle-phonon nuclear model (QPNM) [15] for the case of finite temperature. The set of FT-RPA equations obtained in the framework of this approach will be used in our calculations.

*In nuclei such a phase transition means that in the thermal nuclear characteristics there are two regions with different energy dependences [3].

We consider the model Hamiltonian consisting of the terms describing the motion of nucleons in the average nuclear field, the superfluid pairing interaction and the residual interaction in the form of the separable multipole isoscalar and isovector forces. The form of the average nuclear field is described in calculations by the Woods-Saxon potential $U(r)$. The isoscalar and isovector multipole forces are expressed in terms of the multipole moment operators given by [15]

$$M^{+}_{\lambda\mu} = \frac{(-)^{\lambda-\mu}}{\sqrt{2\lambda+1}} \sum_{jj'} f^{(\lambda)}_{jj'} \{ \frac{1}{2} u_{jj'} [A^{+}(jj';\lambda\mu) + (-)^{\lambda-\mu} A(jj';\lambda-\mu)] + v_{jj'} B(jj';\lambda\mu) + \sum_{m} (-)^{j-m} \delta_{jj'} v^{2}_{j} \}, \tag{1}$$

$$A(jj';\lambda\mu) = \sum_{mm'} (jmj'm'|\lambda\mu) \alpha_{j'm'} \alpha_{jm},$$

$$B(jj';\lambda\mu) = \sum_{mm'} (-)^{j'+m'} (jmj'm'|\lambda\mu) \alpha^{+}_{jm} \alpha_{j'-m'}.$$

In equation (1) α^{+} and α are the creation and annihilation quasiparticle operators; $u_{jj'} \equiv u_j v_{j'} + u_{j'} v_j$, $v_{jj'} \equiv u_j u_{j'} - v_j v_{j'}$ are combinations of Bogolubov's coefficients u and v; $f^{(\lambda)}_{jj'}$ are reduced matrix elements of the single-particle operators generating excitations of the multipolarity λ. We use for the radial dependence of $f^{(\lambda)}_{jj'}$ the derivative $\partial U(r)/\partial r$. As has been noted by Ref. [16] such a radial dependence provides an intensification of the residual interaction in the surface region of the nucleus.

We define the creation and annihilation phonon operators as

$$Q^{+}_{\lambda\mu i} = \frac{1}{2} \sum_{jj'} \{ \psi^{\lambda i}_{jj'} [u_{jj'} A^{+}(jj';\lambda\mu) + v_{jj'} B(jj';\lambda\mu)] - (-)^{\lambda-\mu} \phi^{\lambda i}_{jj'} [u_{jj'} A(jj';\lambda-\mu) + v_{jj'} B^{+}(jj';\lambda-\mu)], \tag{2}$$

$Q_{\lambda\mu i}$ is the Hermitian conjugated operator of $Q^{+}_{\lambda\mu i}$.

Equation (2) shows that unlike the zero temperature phonon operators, defined as usual [15] the phonon operators (2) contain terms with $v_{jj'} B(jj';\lambda\mu)$ and $v_{jj'} B^{+}(jj';\lambda-\mu)$.

To obtain the set of FT-RPA equations for finding the self-energies $\omega_{\lambda i}$ of operators (2) we express the model Hamiltonian in terms of operators (2) and use the commutation relations for operators A and B in equations (1):

$$\langle [A(j_1 j'_1;\lambda_1\mu_1), A^{+}(j_2 j'_2;\lambda_2\mu_2)] \rangle = \delta_{\lambda_1\lambda_2} \delta_{\mu_1\mu_2} (\delta_{j_1 j_2} \delta_{j'_1 j'_2} - (-)^{j'_1+j_1-\lambda_2} \delta_{j'_1 j_2} \delta_{j'_2 j_1})(1 - n_{j_1} - n_{j'_1}), \tag{3}$$

$$\langle [B^{+}(j_1 j'_1;\lambda_1\mu_1), B(j_2 j'_2;\lambda_2\mu_2)] \rangle = \delta_{\lambda_1\lambda_2} \delta_{\mu_1\mu_2} \delta_{j_1 j_2} \delta_{j'_1 j'_2} (n_{j'_1} - n_{j_1}).$$

The brackets $\langle \ldots \rangle$ denote the average value over the grand canonical ensemble, n_j is the occupation probability of the quasiparticle state with the energy ϵ_j : $n_j = [1 + \exp(\epsilon_j/T)]^{-1}$.

The set of FT-RPA equations derived by linearizing the equations of motion for operators (2) is now given by

$$(\kappa^{(\lambda)}_0 + \kappa^{(\lambda)}_1)[X^{\lambda i(n)}_T(\omega_{\lambda i}) + X^{\lambda i(p)}_T(\omega_{\lambda i})] - 4\kappa^{(\lambda)}_0 \kappa^{(\lambda)}_1 X^{\lambda i(n)}_T(\omega_{\lambda i}) X^{\lambda i(p)}_T(\omega_{\lambda i}) = 1,$$

$$X^{\lambda i}_T(\omega_{\lambda i}) = \frac{1}{2\lambda+1} \sum_{jj'} [f^{(\lambda)}_{jj'}]^2 [u^2_{jj'}(1 - n_j - n_{j'}) \frac{\epsilon_j + \epsilon_{j'}}{(\epsilon_j + \epsilon_{j'})^2 - \omega^2_{\lambda i}} - v^2_{jj'}(n_j - n_{j'}) \frac{\epsilon_j - \epsilon_{j'}}{(\epsilon_j - \epsilon_{j'})^2 - \omega^2_{\lambda i}}]. \tag{4}$$

In equations (4) $\kappa^{(\lambda)}_0$ and $\kappa^{(\lambda)}_1$ stand for the isoscalar and isovector constants of multipolarity λ, respectively. The superscripts n or p denote neutrons or protons. The sum is carried out over all quantum numbers $j \equiv (n, \ell, j)$ and $j' \equiv (n, \ell, j')$ of quasiparticles on the average field levels. System (4) coincides exactly with the systems of equations obtained by Ignatyuk [6] with the use of the Green function method or by Sommermann [7] with the use of the method of the density matrix. In the zero temperature limit the occupation probability n_j becomes zero and equations (4) are transformed completely into the set of RPA-equations for cold nuclei [15,17]. The functions $X^{\lambda i(n)}_T(\omega_{\lambda i})$ and $X^{\lambda i(p)}_T(\omega_{\lambda i})$ can be calculated by solving equations (4). The values of the reduced electric transition probabilities $B_T(E\lambda\uparrow, \omega_{\lambda i})$ are calculated in the long-wave length approximation as [15]

$$B_T(E\lambda\uparrow, \omega_{\lambda i}) = \left| \frac{e_{eff} X^{\lambda i(n)}_T(\omega_{\lambda i})}{\sqrt{2\mathcal{Y}_n(\lambda i)}} + \frac{(1 + e_{eff}) X^{\lambda i(p)}_T(\omega_{\lambda i})}{\sqrt{2\mathcal{Y}_p(\lambda i)}} \right|^2, \tag{5}$$

where

$$\mathcal{Y}_{n,p}(\lambda i) = Y^{\lambda i}_{n,p} + Y^{\lambda i}_{p,n}\left\{\frac{(\kappa^{(\lambda)}_0 - \kappa^{(\lambda)}_1)X^{\lambda i(n,p)}_T(\omega_{\lambda i})/(2\lambda+1)}{1-[(\kappa^{(\lambda)}_0 + \kappa^{(\lambda)}_1)X^{\lambda i(p,n)}_T(\omega_{\lambda i})/(2\lambda+1)]}\right\},$$

$$Y^{\lambda i}_{n,p} = \frac{1}{2}\partial X^{\lambda i(n,p)}_T/\partial\omega|_{\omega=\omega_{\lambda i}} .$$

The electric effective charge e_{eff} is given for E1-transitions as $e^{(1)}_p = N/A$, $e^{(1)}_n = -Z/A$, and for E2-transitions as $e^{(2)}_p = 1-(A+N)/A^2$, $e^{(2)}_n = Z/A^2$.

In the FT-RPA moments are defined as

$$m_k(T) = \sum_{\lambda i}\omega^k_{\lambda i}B_T(E\lambda\uparrow,\omega_{\lambda i}), \tag{6}$$

where the summation runs over all one-phonon states with energies $\omega_{\lambda i}$. The centroid $\bar{E}$ of the strength distribution and the Landau splitting σ are defined, respectively, as

$$\bar{E}_T = m_1(T)/m_0(T), \quad \sigma = \{[m_2(T)/m_0(T)] - \bar{E}^2_T\}^{1/2}. \tag{7}$$

Detailed studies of the temperature dependence for the moments have been given in Refs.[9-12,18]. They have shown that the first moment (the EWSR) is temperature independent for dipole states and has a weak temperature dependence for quadrupole states. We made our calculations based on equations (4-7). The single-particle energies E_j are defined using the set of parameters for the Woods-Saxon potential chosen at zero temperature by Cherpunov[19] and Takeuchi and Moldauer[20]. This single-particle energy spectrum is used at finite temperature since its temperature dependence is very smooth and weak up to T = 6 MeV, as has been shown in Ref.[21]. The superfluid pairing forces are taken into account by solving the BCS-equations at finite temperature[6]

$$\Delta_T = G\sum_j(j+\frac{1}{2})u_jv_j(1-2n_j).$$

The critical temperature T_{crit}, for which the collapse of the pairing gap takes place, is given by[6,14] $T_{crit} = 0.567\Delta(T=0)$. The excitation energy E^* of the equilibrium state for each temperature has the form $E^* = \mathcal{E}(T) - \mathcal{E}(T=0)$, where $\mathcal{E}(T)$ is the energy of the system at temperature T

$$\mathcal{E}(T) = \sum_j(2j+1)E_j\left[1-\frac{E_j-\lambda}{\epsilon_j}(1-2n_j)\right] - \Delta^2_T/G,$$

λ is the chemical potential. The quasiparticle energy ϵ_j is calculated through the single particle energy E_j and the gap Δ_T as follows $\epsilon_j = [(E_j-\lambda)^2 + \Delta^2_T]^{1/2}$.

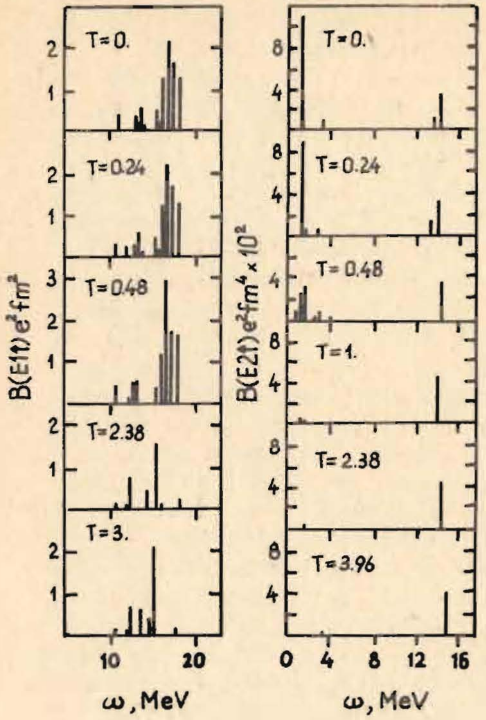

Fig.1. RPA strength distributions for IV dipole and IS quadrupole modes at different temperatures in ^{58}Ni.

We choose the constant $\kappa^{(\lambda)}_0$ and $\kappa^{(\lambda)}_1$ for the dipole and quadrupole modes as follows. At zero temperatures the isoscalar energy $\omega^{(1)}_0$ for dipole modes is put to be equal to zero to exclude the "spurious" states. The isovector dipole $\kappa^{(1)}_1$ constant is defined to reproduce the empirical location of the isovector dipole resonance. After that the ratio $\kappa^{(1)}_1/\kappa^{(1)}_0$ is used for finding the isoscalar and isovector quadrupole $\kappa^{(2)}_{0,1}$ constants so as to reproduce the lowest 2^+_1-energy, defined empirically for each nucleus at zero temperature by Sakai[22]. Calculations performed by Blokhin and Ignatyuk[23] have shown that the multipole constants $\kappa^{(\lambda)}$ have a very weak temperature dependence. Therefore, we use the zero-temperature constants $\kappa^{(\lambda)}_{0,1}$ all through our calculations.

The strength distributions of the isovector (IV) dipole E1 and isoscalar (IS) quadrupole excitations as the distributions of the reduced transition probabilities $B(E1\uparrow)$ and $B(E2\uparrow)$ at several temperatures T in ^{58}Ni are shown in fig.1. The RPA result at T = 0 MeV gives a collective E1 resonance near $\omega_{\lambda i}$ = 17 MeV exhausting 80% of the EWSR and a collective E2 resonance near $\omega_{\lambda i}$ = 14 MeV exhausting 72% of the IS EWSR in the energy interval $10 \le \omega_{\lambda i} \le 25$ MeV. Empirical evidence has been reported at around $\omega_{\lambda i}$ = 17.3 MeV having (94±10)% of the EWSR value in the energy interval $10 \le \omega_{\lambda i} \le 30$ MeV for the giant dipole resonance (GDR)[24,25]. The giant quadrupole resonance (GDR) has been extracted at around $\omega_{\lambda i}$ = 14 MeV by Kocher and Auble[24] and at around $\omega_{\lambda i}$ = 16 MeV having (65±10)% of the IS EWSR value in the energy interval $10 \le \omega_{\lambda i} \le 30$ MeV by Pitthan et al.[25]. While at zero temperature the quadrupole strength of the lowest 2^+_1-states is mainly concentrated on a single level at 1.4545 MeV[22], at finite temperature it is spread over several states with smaller B(E2↑)values. As has been shown by Ignatyuk [6] and Sommermann [7] at finite temperature the poles of the (p-h) excitations at energies $\epsilon_j + \epsilon_{j'}$ remain essentially in their original location, but simultaneously the poles of (p-p) and (h-h) excitations appear at energies $\epsilon_j - \epsilon_{j'}$ (see equation (4)). The additional poles are caused by the terms containing $v_{jj'}$ $B(jj';\lambda\mu)$ and $v_{jj'}B^+(jj';\lambda-\mu)$ in eq.(1). Our calculations show that the new poles at energies $\epsilon_j - \epsilon_{j'}$ correspond to quite noncollective (p-p) and (h-h) configurations with very weak B(Eλ↑)-values. Thus, for the quadrupole modes, for example, at T ~ 0.3 MeV the B(E2↑)-values due to the states appearing between the new poles $\epsilon_j - \epsilon_{j'}$ at energies <1 MeV are smaller than 10^{-4} % of the B(E2↑)-value for 2_1^+-state in ^{58}Ni defined at T = 0 MeV. However, the new configurations lead to a less transition strength and to a spreading of the giant resonances at higher temperature. The location of the giant resonances is shifted to the lower energy side for the IV GDR and to the higher one for the IS GDR (fig.1). This is in agreement with the general feature of the results obtained in Refs. [8-11,13,26].

To investigate the behaviour of the giant resonance energy and their widths in the dependence on temperature including the critical region due to the phase transition from the superfluid state to the normal one we show in the Table and figure 2a the centroids $\bar{E}$ of the main peaks of IV dipole and IS quadrupole resonances in the region near T_{crit} = 0.79 MeV. As has been shown by Landau and Lifshitz [14] and for nuclei by Ignatyuk [6], this phase transition is a second order one because of the presence of a break in the behaviour of the excitation energy E* at this critical point, corresponding to a finite discontinuity in the first order derivative for the entropy (see fig.3). In the region $T \geq T_{crit}$ the gap Δ is zero. Our results show that while the $\bar{E}$ behaviour of giant resonances has a smooth and weak temperature dependence at T far from T_{crit}, it undergoes a bending in the critical region and a break at $T = T_{crit}$ (fig.2a). The jump of the $\bar{E}$ behaviour at T near T_{crit} is somewhat larger for the IV dipole resonance as compared with the one for the IS quadrupole resonance. The occurrence of this bending in the behaviour of giant resonances is undoubtedly an important property for superfluid at zero temperature nuclei. In closed-shell nuclei, for which the pairing gap Δ(T = 0) is zero, there is no such a discontinuity and the temperature dependence of giant resonance energies gives smooth curves, as has been shown in Refs. [7,8,10]. This discontinuity is absent also in the calculations, where the superfluid pairing forces are not taken into account or are temperature independent. This is confirmed by the dashed lines in figs.2a and 3 for the $\bar{E}$-behaviour of giant resonances and the excitation energy E*, which are obtained

Table

RPA results for IV dipole (E1) and IS quadrupole (E2) states in ^{58}Ni at finite temperature. a - Calculations with the superfluid gap Δ_T, obtained by solving the finite temperature BCS-equation. b - Calculations with $\Delta_T = \Delta(T = 0)$. All the values are given in MeV

T		0.24	0.48	0.55	0.63	0.71	0.74	0.79	1.00	2.38
E*	a	0.28	4.92	7.03	8.96	10.50	10.99	11.60	12.49	37.16
	b	0.22	4.34	6.23	8.00	9.59	10.13	10.95	13.72	40.32
$\bar{E}_1$	a	16.88	16.88	16.82	16.73	16.68	16.57	15.65	15.64	15.50
	b	16.88	16.88	16.84	16.83	16.81	16.80	16.79	16.75	16.43
σ_1	a	1.56	1.65	1.32	1.54	1.25	1.24	1.58	1.67	1.99
$\bar{E}_2$	a	14.04	14.02	13.99	13.95	13.90	13.85	13.82	13.85	14.13
	b	14.04	14.04	14.05	14.07	14.07	14.08	14.08	14.10	14.30
σ_2	a	0.58	0.54	0.64	0.69	0.89	1.29	0.71	0.61	0.23

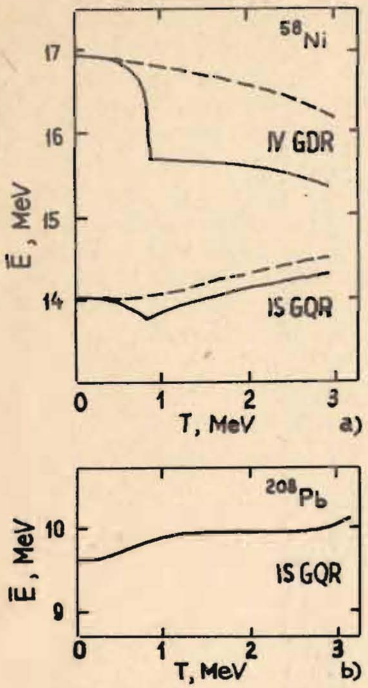

Fig.2. a - Centroids $\bar{E}$ *of the IV dipole and IS quadrupole resonances as functions of temperature in* ^{58}Ni. *Dashed lines correspond to the case* $\Delta_T = \Delta(T=0)$ *b - Centroid* $\bar{E}$ *of the IS quadrupole resonances as a function of temperature in* ^{208}Pb.

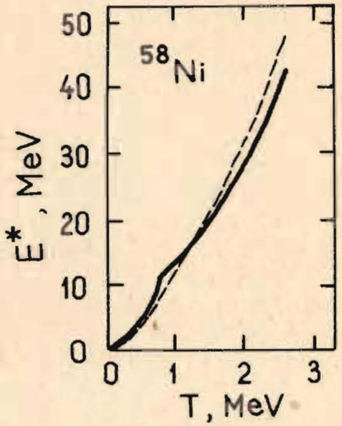

Fig.3. Excitation energy E* *as a function of temperature, obtained in* ^{58}Ni *by solving the temperature dependent gap equation (solid line), and with* $\Delta_T = \Delta(T=0)$ *(dashed line).*

with $\Delta_T = \Delta(T=0)$, and by fig.2b which presents the $\bar{E}$ behaviour of the IS quadrupole resonances in ^{208}Pb (cf.also Refs.[8,10]).Ring et al.[27] have shown that in rotating nuclei at finite temperature the pairing forces can play also an important role in the behaviour of GDR. We note that in rotating nuclei there is another kind of phase transition connected with the nuclear deformations. Thus, the features of nuclear deformations produced by the alignment of individual particles or pairs and the first order phase transition in rotating nuclei have been investigated in detail in Ref.[28].

The analysis of the experimental results, recently reviewed by Snover[29], show a large broadening of the GDR width when the temperature increases as well as a downward shift by 1-1.5 MeV. Our calculations, which give a downward shift of the centroid of the IV dipole resonance and a spreading of collective states at higher temperatures are in agreement with these observations. The values of the Landau splitting in our calculations tend to increase for IV dipole resonance and to decrease for IS quadrupole resonance in the region $T > T_{crit}$ as the temperature T increases (the Table). However, for a correct study of the temperature dependence of the giant resonance width one must consider both escape $\Gamma^{\uparrow}$ and spreading $\Gamma^{\downarrow}$ widths. The first is caused by taking into account the continuum. The broadening of the giant resonance width at finite temperature appearing as a consequence of the effects of the continuum has been investigated recently by Sagawa and Bertsch [11] in the framework of the self-consistent temperature response function. The spreading width $\Gamma^{\downarrow}$ is caused by the interaction with more complicated configurations. In the QPNM the interaction with 2p-2h configurations at zero temperature is studied in detail by taking into account two-phonon components in the wave function excitation operator[15]. In this way a description of the giant resonance spreading width has been given in good agreement with the experimental data[30]. It is highly desirable to develop this method for describing the giant resonance spreading width at finite temperature. It is a goal for future studies.

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Нгуен Динь Данг E4-85-194

Влияние спаривания сверхпроводящего типа на коллективные состояния в приближении хаотических фаз при конечной температуре

В рамках квазичастично-фононной модели ядра получена система уравнений в приближении хаотических фаз при конечной температуре. Рассчитаны изовекторные дипольные и изоскалярные квадрупольные моды в ^{58}Ni. Найдено, что кривая зависимости центроидов гигантских резонансов от температуры претерпевает изгиб в области критической температуры, где разрушается спаривание сверхпроводящего типа.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

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Nguyen Dinh Dang E4-85-194

Influence of the Superfluid Pairing Interaction on the Collective States in the Finite Temperature Random Phase Approximation

A set of finite temperature RPA-equations is obtained within the framework of the quasiparticle-phonon nuclear model using the separable multipole interaction. The calculations have been performed for the isovector dipole and isoscalar quadrupole modes in ^{58}Ni. We find that the centroids of the giant resonances undergo a bending in the critical temperature region where the collapse of the pairing gap takes place.

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

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