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Research Article

Application of Correlated States of Interacting Particles in Nonstationary and Periodical Modulated LENR Systems

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Abstract

In the report the universal mechanism of optimization of low energy nuclear reactions (LENR) on the basis of coherent correlated states of interacting particles at different kinds of nonstationary periodical action to the system is considered. We have considered the peculiarities and investigated the efficiency of the creation of a correlated state under a periodic action on a harmonic oscillator. This method is shown to lead to rapid formation of a strongly correlated particle state that provides an almost complete clearing of the potential barrier even for a narrow range of oscillator frequency variations. Several successful low-energy fusion experiments based on usage of correlated states of interacting particles are discussed.

Keywords: Coherent correlated states, LENR, Transparency of Coulomb potential barrier

1. Introduction

The problem of overcoming the Coulomb potential barrier during the interaction of charged particles is the key one in modern atomic and nuclear physics. This problem is of particular importance in controlled nuclear fusion, for which the action of the Coulomb barrier leads to a very low tunneling probability for low energy particles.

This fundamental constraint leads to extremely big problems related to the necessity of very rapid heating of a target composed of hydrogen isotopes (deuterium, D and tritium, T) to a state of thermonuclear plasma with a temperature of about 10 keV and its additional compression by a factor of $k \geq 10^3$ (in the case of "inertial" nuclear fusion with a confinement time of 10^{-9} s based on an initial target with an atomic concentration of 10^{23} cm^{-3}) as well as similar heating of a low density target to the same temperature with its subsequent confinement for a long time (for TOKAMAK or STELLARATOR systems).

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Extremely big obstacles that should be overcome for a successful solution of such global problems in future are well known. In addition, the studies of controlled d + d and d + t fusion carried out for 60 years have shown that using the "thermonuclear" way of overcoming the Coulomb barrier makes any attempt to use the controlled thermonuclear fusion reactions based on isotopes heavier than deuterium or tritium under terrestrial conditions virtually unrealistic [1,2].

It is well known that the d + d and d + t fusion reactions are not optimal due to an intense flux of fast neutrons accompanying these reactions. This leads to induced radioactivity and critical radiation damage of the reactor vessel (the first wall problem). In this respect, for example, fusion based on the ecologically pure reaction

$$\mathrm{B}^{11} + \mathrm{p} = 3\mathrm{He}^4 + 8.7\ \mathrm{MeV} \tag{1}$$

involving B^{11} and p nuclei is more optimal: it does not produce neutrons and does not induce activity.

In these circumstances, the topicality of new approaches that can provide an increase in the transparency of potential barriers through the application of quantum processes determining the interaction of low energy particles with these barriers without any compulsory use of a high temperature plasma under conditions of either long confinement or super-strong compression is obvious. In [1–3], we have shown that the use of coherent correlated states can lead to such effect.

A specific quantitative analysis of the condition for the formation of a strongly correlated state has been performed in [1–4] basing on one particular and specially chosen deformation regime of an oscillator with a particle located in its parabolic field irrespective of the possibility of its practical realization. Below, we consider in detail the possible methods of the formation of such states, determine the conditions of their realization, and show that there exist many different and simple (compared to [3]) realization scenarios close to the conditions of a real experiment.

2. The General Foundations of Formation of Coherent Correlated States of Particles

The Heisenberg uncertainty relation for the coordinate and momentum,

$$\sigma_q \sigma_p \ \geq \ \hbar^2/4 \tag{2a}$$

and its generalization

$$\sigma_A \sigma_B \ \geq \ |\langle[\widehat{A}\widehat{B}]\rangle|^2/4, \quad \sigma_C = \langle(\Delta\widehat{C})^2\rangle \equiv (\delta C)^2, \ \ \Delta\widehat{C} = \widehat{C} - \langle C\rangle \tag{2b}$$

made in 1929 by Robertson for arbitrary dynamical variables A and B are base relations of quantum mechanics. In modern interpretation these relations correspond to uncorrelated states.

In 1930, Schrodinger [5] and Robertson [6] generalized relation (2b) and derived a more universal inequality called the Schrodinger–Robertson uncertainty relation

$$\sigma_A \sigma_B \ \geq \ |\langle[\widehat{A}\widehat{B}]\rangle|^2/4(1 - r^2), \tag{3a}$$

$$r = \sigma_{AB}/\sqrt{\sigma_A \sigma_B},$$

$$\begin{aligned}\sigma_{AB} &= \langle\{\Delta\widehat{A},\, \Delta\widehat{B}\}\rangle/2\\ &\equiv \langle((\widehat{A}-\langle A\rangle)\,(\widehat{B}-\langle B\rangle)+(\widehat{B}-\langle B\rangle)\,(\widehat{A}-\langle A\rangle))\rangle/2\\ &= (\langle\widehat{A}\widehat{B}+\widehat{B}\widehat{A}\rangle)/2-\langle A\rangle\,\langle B\rangle,\end{aligned}$$

where r is the correlation coefficient between quantities $A$and B with $|r| \le 1$, σ_{AB} is the mutual variance of A and B corresponding to the mean value of the anticommutator of the error operators $\Delta\widehat{K} = \widehat{K} - \langle K\rangle$.

The Schrodinger–Robertson uncertainty relation (3a) is an obvious generalization of the Heisenberg–Robertson uncertainty relation (2b) for correlated states and is reduced to it at $r = 0$.

In [1–4,7,8] it was shown that for a model system including a particle with coordinate $q(t)$ and momentum $p(t)$ in the field of a nonstationary harmonic oscillator

$$V(q,t) = \frac{mq(t)^2\omega^2(t)}{2}, \tag{3b}$$

a decrease in the particle oscillation frequency $\omega(t)$, leads to an increase in the correlation coefficient $|r(t)|$, and a change of the uncertainty relation,

$$r(t) = \langle q\widehat{p}_q + \widehat{p}_q q\rangle/2\delta q\delta p_q,$$

$$\delta q \equiv \sqrt{\langle q^2\rangle}, \qquad \delta p_q \equiv \sqrt{\langle p_q^2\rangle},$$

$$\delta q\ \delta p_q \ge \hbar/2\sqrt{1-r^2}. \tag{3c}$$

Formally, the change in the correlation coefficient in the uncertainty relation can be taken into account by the formal substitution

$$\hbar \to \hbar^* \equiv \hbar/\sqrt{1-r^2}. \tag{4}$$

In the absence of a correlation the uncertainty relation takes the form of the "standard" Heisenberg relation (2a).

When a strongly correlated particle state with $|r| \to 1$ is formed, the product of the variances of the coordinate $\langle q^2\rangle$ and momentum $\langle p_q^2\rangle$ increase indefinitely. This leads to the possibility of a much more efficient particle penetration into the sub-barrier region $V(q)$ than that for the same particle in an uncorrelated state.

It was shown in [1–4] that the very low barrier transparency (tunneling probability) for the initial uncorrelated state,

$$D_0 \equiv D_{r=0} = \exp\{-W(E)\} \ll 1, \tag{5}$$

$$W(E) = \frac{2}{\hbar}\int_R^{R+L(E)} |p(q)|\,\mathrm{d}q \equiv \frac{2}{\hbar}\,\langle|p(q)|\rangle\, L(E),$$

$$|p(q)| = \sqrt{2M}\left\langle\sqrt{V(q)-E}\right\rangle$$

that corresponds to the conditions $E \ll V_{\max}$, $W(E) \gg 1$ for the formation of a strongly correlated superposition particle state can increase to a very large value, $D_{|r|\to 1} \to 1$ at the same low energy $E \ll V_{\max}$.

In (5) R is the nuclear radius, L is the "barrier width" and M is the reduced particle mass.

In a very simplified form, this effect can be taken into account by the formal (not quite correct) substitutions

$$W_{r=0}(E) \to W_{r\neq 0}(E,\hbar) \equiv W_{r=0}(E,\hbar^*) = W_{r=0}(E,\hbar)\sqrt{1-r^2} \tag{6}$$

in (5) and

$$D_{r\neq 0} \approx \exp\left\{-\frac{2\sqrt{1-r^2}}{\hbar}\int\limits_R^{R+L(E)} |p(q)|\,\mathrm{d}q\right\}, \tag{7}$$

from which it follows that

$$D_{|r|\to 1} \to 1 \ \text{ even if } \ E \ll V_{\max} \ \text{ and } \ W_{r=0}(E) \gg 1. \tag{8}$$

In this case, the potential barrier transparency increases by a factor of

$$D_0^{\sqrt{1-r^2}}/D_0 = 1/D_0^{1-\sqrt{1-r^2}} \gg 1, \tag{9}$$

which is close in order of magnitude to the result of exact barrier clearing calculations using rigorous quantum-mechanical methods [1,3,4]. Although these estimates with the substitution $\hbar \to \hbar^*$ are not quite correct (they are made just for illustration of order of the effect) and must be justified every time, they clearly demonstrate a high efficiency of the use of coherent correlated states in solving of applied tunneling related problems in the case of a high potential barrier and a low particle energy.

The physical causes of the huge increase in barrier transparency for a particle in coherent correlated superposition state were discussed in our works [1,3,4]. The physical reason for the increase of the probability of tunneling effect is related to the fact that the formation of a coherent correlated state leads to the cophasing and coherent summation of all fluctuations of the momentum for various eigenstates forming the superpositional correlated states. This leads to great dispersion of the momentum correlated state, very great fluctuations of kinetic energy of the particle in the potential well and increasing of potential barrier penetrability.

A coherent correlated state can be formed in various quantum systems. The most easy way to form such state is when the particle is in a nonstationary parabolic potential well (3b).

The formation mechanism in such system was considered in [1–4,7,8].

There are two different methods for analyzing and optimizing the process of formation of a correlated state.

The first method [1] is that a prespecified form of the specific dependence $r(t)$ that must simultaneously satisfy condition of adiabaticity

$$\left|\frac{\mathrm{d}r(t)}{\mathrm{d}t}\right| \ \le \ \omega_0\sqrt{1-r(t)^2}\Big/\left|\frac{r(t)}{|r(t)|} + 2\omega_0\int\limits_0^t \frac{r(t')\,\mathrm{d}t'}{\sqrt{1-(r(t'))^2}}\right| \tag{10}$$

and the optimal condition $|r| \to 1$ should be used to find the optimal condition for the change in

$$\omega(t) = A\sqrt{\frac{1}{(g(t))^2(1-r^2)}\left\{1 - \frac{g(t)}{\omega_0\sqrt{1-r^2}}\frac{\mathrm{d}r}{\mathrm{d}t}\right\}}. \tag{11}$$

Although this method seems to be logically correct, it is not optimal. In most cases the prespecified "convenient" form of the dependence $r(t)$ is inconsistent with the adiabaticity condition (10) or leads to an expression for $\omega(t)$ that is very difficult to realize in an experiment.

The second more optimal method [2–4] consists in solving the inverse problem-finding $r(t)$ for a specified law of change in the oscillator frequency $\omega(t)$ that is consistent with the possibilities of a real experiment.

In this case, the problem of ensuring the adiabaticity condition is solved automatically by using the real dependence $\omega(t)$. A parametric form of the solution $r(\omega(t))$ inverse to (11) and equivalent to it that is more convenient for calculations can be used to find such solutions.

This solution can be obtained by analyzing the complex equation of motion for a classical oscillator with a variable frequency that in dimensionless form is

$$\frac{\mathrm{d}^2\varepsilon}{\mathrm{d}t^2} + \omega^2(t)\varepsilon = 0, \qquad \varepsilon(0) = 1, \qquad \left.\frac{\mathrm{d}\varepsilon}{\mathrm{d}t}\right|_0 = i, \qquad \omega(0) = \omega_0, \tag{12}$$

where ε is the complex amplitude of the harmonic operator normalized to $x_0 = \sqrt{\hbar/M\omega_0}$.

The correlation coefficient is defined by the expression [1–3,6,7]

$$r = Re\left\{\varepsilon^*\frac{\mathrm{d}\varepsilon}{\mathrm{d}t}\right\} / \left|\varepsilon^*\frac{\mathrm{d}\varepsilon}{\mathrm{d}t}\right|, \qquad r^2 = 1 - \omega_0^2/\left|\varepsilon^*\frac{\mathrm{d}\varepsilon}{\mathrm{d}t}\right|^2. \tag{13}$$

The general solution of (12) is

$$\varepsilon(t) = \mathrm{e}^{\varphi(t)}, \qquad \varphi(t) = \alpha(t) + i\beta(t). \tag{14}$$

Substituting (14) into (12) and (13), using the initial conditions

$$\varphi(0) = \alpha(0) = \beta(0) = 0,$$

$$\left.\frac{\mathrm{d}\varphi}{\mathrm{d}t}\right|_0 = i, \qquad \left.\frac{\mathrm{d}\alpha}{\mathrm{d}t}\right|_0 = 0, \qquad \left.\frac{\mathrm{d}\beta}{\mathrm{d}t}\right|_0 = 1 \tag{15}$$

that follow from (12), and separating the real and imaginary parts of the derived equation, we find

$$\frac{\mathrm{d}^2\alpha}{\mathrm{d}t^2} + \left(\frac{\mathrm{d}\alpha}{\mathrm{d}t}\right)^2 - \exp(-4\alpha) = -\omega^2(t), \tag{16a}$$

$$\beta(t) = \int_0^t \exp\{-2\alpha(t')\}\,\mathrm{d}t', \tag{16b}$$

$$|r| = \sqrt{\frac{\left(\frac{d\alpha}{dt}\right)^2 \exp(4\alpha)}{\left\{1+\left(\frac{d\alpha}{dt}\right)^2 \exp(4\alpha)\right\}}}. \tag{16c}$$

The system of Eqs. (16a)–(**??**) is equivalent to Eq. (11), but it is more convenient to analyze and allows one initially to find the exponent of the oscillation amplitude for the oscillator $\alpha(t)$ from Eq. (16a) based on a specified law of change in frequency $\omega(t)$ and subsequently to find $r(t)$ from Eq. (**??**) based on $\alpha(t)$.

Note that our analysis disregards the influence of the damping of the oscillations of a nonstationary harmonic oscillator on the formation of a correlated state of the quantum system. This damping can be taken into account by analyzing a more general equation of oscillator motion

$$\frac{d^2\varepsilon}{dt^2} + 2\gamma\frac{d\varepsilon}{dt} + \omega^2(t)\varepsilon = 0, \qquad \varepsilon(0) = 1, \qquad \left.\frac{d\varepsilon}{dt}\right|_0 = i, \qquad \omega(0) = \omega_0 \tag{17a}$$

and the equations following from it,

$$\frac{d^2\alpha}{dt^2} + \left(\frac{d\alpha}{dt}\right)^2 + 2\gamma\frac{d\alpha}{dt} - \exp(-4\alpha - 4\gamma t) = -\omega^2(t), \tag{17b}$$

$$\beta(t) = \int_0^t \exp\{-2\alpha(t') - 2\gamma t'\}\,dt'. \tag{17c}$$

The problem of influence of both damping and presence of additional fluctuation force was solved also [4] and will be discussed later.

3. Formation of Correlated States of Interacting Particles under Nonstationary Periodical Action

Using solutions (16a)–(16c), we will determine the formation dynamics of correlated states for different regimes of change in frequency $\omega(t)$.

In our works [1,3] the method of formation of coherent correlated states of particle at monotonic decrease in the frequency $\omega(t)$ of a nonstationary harmonic oscillator (e.g. $\omega(t) = \omega_0 \exp(-t/T)$) was discussed.

A more interesting and realizable situation takes place for a harmonic law of change in $\omega(t)$ in the case of a full-scale (maximally possible) change of the oscillator frequency,

$$\omega(t) = \omega_0\,|\cos\Omega t| \tag{18}$$

or a change of this frequency in a limited range,

$$\omega(t) = \omega_0\,(1 + g_\Omega \cos\Omega t), \tag{19}$$

where $|g_\Omega| < 1$ is the modulation depth.

This regime can be provided, for example, at a constant depth of the potential well $V_{\max}$ in which the particle is located and for a periodic change in its width in the interval

$$L_0/(1+|g_\Omega|) \le L \le L_0/(1-|g_\Omega|), \;\; L_0 = \sqrt{8V_{\max}\Big/M\omega_0^2}. \tag{20}$$

The efficiency of excitation of the correlated states greatly depends on a ratio of frequencies ω_0 and Ω.

Figure 1 presents the time dependences of the correlation coefficient $r(t)$ for a periodic and limited change in the oscillator frequency (19) at $\Omega = \omega_0, 2\omega_0$ and at various frequency modulation depths g_Ω.

It follows from these results that a completely correlated state is formed not only for a monotonic–asymptotic [1,3] or periodic full-scale (in the range $0 \le \omega(t) \le \omega_0$) change in the oscillator frequency, but also for its change in the limited range

$$\omega_0\,(1-g_\Omega) \le \omega(t) \le \omega_0\,(1+g_\Omega) \tag{21}$$

and even for$|g_\Omega| \ll 1$ [3,4].

The duration of formation of correlated state decreases with the increase of frequency modulation depth and reaches its minimum for $|g_\Omega| \to 1$. For example, for the case, presented on Fig. 1(d) ($g_\Omega = 0.1,\;\; \Omega = 2\omega_0$), we have $|r|_{\max} = 0.999998$ at $\omega_0 t = 500$. For such value of $|r|_{\max}$ the probability of tunneling effect for two possible reactions

$$\mathrm{d + d \to He^3 + n, \qquad t + p}, \tag{22}$$

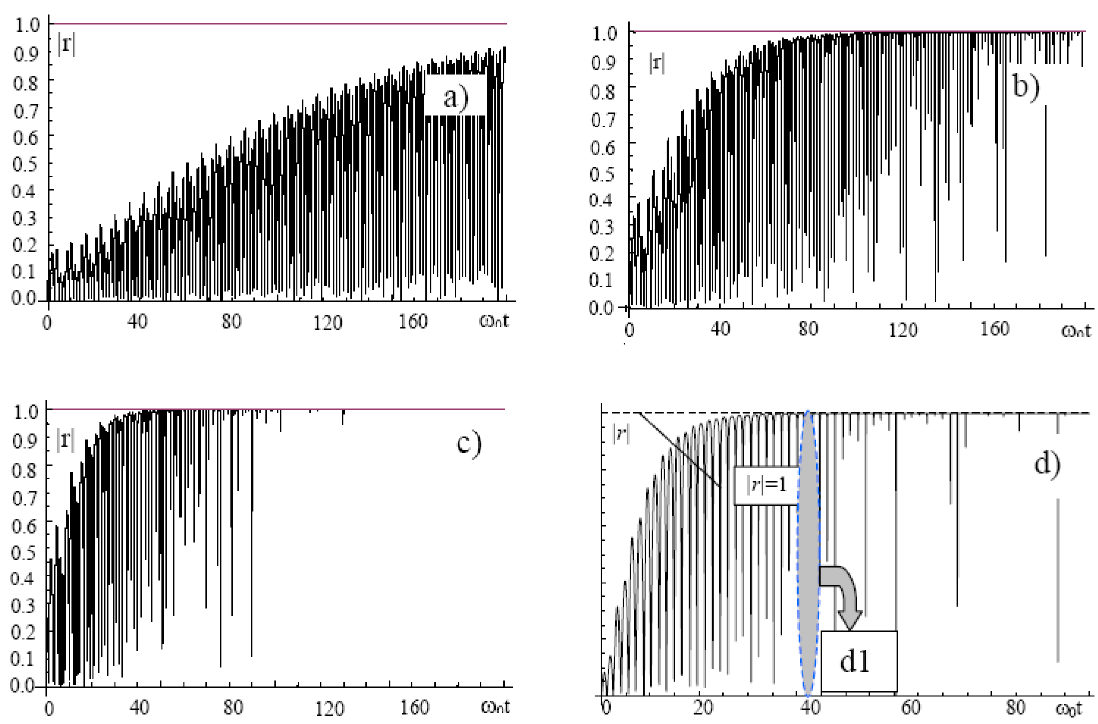

Figure 1(a–d)

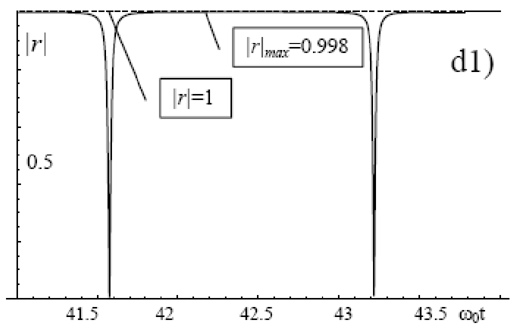

Figure 1. Time dependences of the correlation coefficient $r$for a limited change in the oscillator frequency $\omega(t) = \omega_0\,(1 + g_\Omega \cos\Omega t)$ for various frequency modulation depths g_Ω and various frequency modulation: (a) $g_\Omega = 0.1,\ \ \Omega = \omega_0$; (b) $g_\Omega = 0.2,\ \ \Omega = \omega_0$;(c) $g_\Omega = 0.3,\ \Omega = \omega_0$; (d) $g_\Omega = 0.1,\ \Omega = 2\omega_0$, (d1) increased fragment of (1d).

$$\mathrm{Pd}^A + \mathrm{d} \to \mathrm{Ag}^{A+2} \to \ldots \tag{23}$$

at room temperature increases from $D_{r=0} \approx 10^{-100}$ (for non-correlated state of interacting deuterons in reaction (22)) to $D_{r=0.999998} \approx 0.8$ (for correlated state of d) and from $D_{r=0} \approx 10^{-4600}$ (for non-correlated state of interacting particles d and Pd^A in reaction (23)) to $D_{r=0.999998} \approx 10^{-8}$ in potential well made of nearest Pd^{A+} ions)!

From the detailed analysis follows that the process of correlated states formation at the action of limited periodic modulation $\omega(t) = \omega_0\,(1 + g_\Omega \cos\Omega t)$ is possible only at any of two conditions:

$$\Omega = \omega_0 \tag{24a}$$

or Ω is close to $2\omega_0$ and lies inside the interval P

$$(2 - g_\Omega)\omega_0 \le \Omega \le (2 + g_\Omega)\omega_0. \tag{24b}$$

The results of calculation of averaged correlation coefficient

$$\langle |r(t)| \rangle_t \equiv \frac{1}{\Delta t} \int\limits_{t_0-\Delta t/2}^{t_0-\Delta t/2} |r(t)|\,\mathrm{d}t \tag{25}$$

are presented on Fig.2a,b for $\Delta t = 10^3/\omega_0$ and different values of modulation depths $g_\Omega = 0.1 \quad (t_0 = 1500/\omega_0)$, and $g_\Omega = 0.01 \quad (t_0 = 10^4/\omega_0)$.

In Fig. 2(b), the results of calculation the averaged coefficient of barrier transparency

$$\langle\langle D(t,\Omega)\rangle_t\rangle = \frac{1}{\sqrt{\pi}\,\delta\Omega} \int \left\{ \frac{1}{\Delta t} \int\limits_{\mathrm{t}_0-\Delta t/2}^{\mathrm{t}_0-\Delta t/2} D(t,\Omega')\,dt \right\} \mathrm{e}^{-\frac{(\Omega-\Omega')^2}{(\delta\Omega)^2}}\,\mathrm{d}\Omega' \tag{26}$$

at non-monochromatic periodic modulation with $\Omega/\delta\Omega = 10$ are presented also

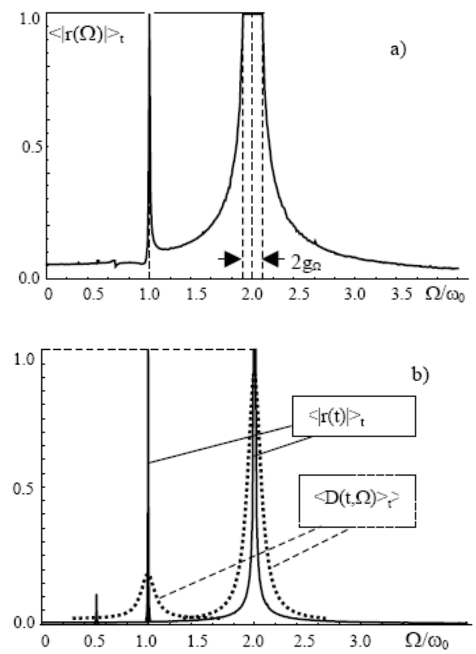

Figure 2. Dependences of averaged correlation coefficient $\langle|r(t,\Omega)|\rangle_t$ and normalized averaged coefficient of barrier transparency on frequency of modulation Ω at $|g_\Omega| = 0.1$ (a) and $|g_\Omega| = 0.01$ (b).

So the essential suppression of action of potential barrier on the effectiveness of nuclear reaction with the participation of charged particles at limited periodic modulation of the oscillator frequency (19) is possible only for two frequencies (24a) and (24b).

From this result very important statement follows: in any experiments with the use of external periodic modulation with limited frequency interval, two resonances of energy release on frequencies $\Omega = \omega_0$ or $|\Omega - 2\omega_0| \leq g_\Omega\omega_0$ should be observed!

This statement are in good correlation with *"Terahertz"* laser experiments of Letts, Hagelstein and Cravens [9,10] on the stimulation of nuclear reaction at joint action of two laser beams with variable beat frequency $\Omega = 3 - 24$ THz on the cathode surface during the electrolysis in PdD system with the presence of heavy water D_2O.

Figure 3 shows the experimental frequency dependencies of thermal energy release [10] in these experiments.

Formation of correlated states in this system is connected with the direct or indirect (by plasmon excitation or phonon mode modulation) action of electromagnetic radiation with frequencies Ω on optical phonon modes $\omega_0^{(k)}$ of deuterons in PdD compound.

Four main resonances of energy release $\Omega_1 \approx 7.8 - 8.2$ THz, $\Omega_2 \approx 10.2 - 10.8$ THz, $\Omega_3 \approx 15.2 - 15.6$ THz and $\Omega_4 \approx 20.2 - 20.8$ THz in Fig. 3 are the result of averaging of about 30 experiments and subsequent statistical processing of experimental data. In work [9] the second resonance with frequency $\Omega_2 \approx 10.2 - 10.8$ THz was not separated.

Comparison of frequencies of all four resonances shows that the ratios between these frequencies are $\Omega_3 \approx 2\Omega_1$ and $\Omega_4 \approx 2\Omega_2$ with good accuracy. By the way, from the given experiments follows that the amplitude of high-frequency maxima in each of these pairs (accordingly Ω_3 and Ω_4) greatly excided the amplitudes of the maxima corresponding to the "basic" frequencies Ω_1 and Ω_2. Such relation directly follows from comparison of Figs. 2(b) and 3. These experimental results completely correspond to theoretical model of the coherent correlated states that was discussed above.

This model also allows to explain the presence of a resonance of nuclear reactions (**??**) and (23) on frequency $\Omega_4 \approx 20.2 - 20.8$ THz (at action of beat frequency Ω_4) . It is known that in the region $\omega_0^{(k)} > 16$ THz there is no optical phonon modes for PdD compound (see analysis in [9]). So, the resonance of nuclear reactions (**??**) and (23) at action of beat frequency Ω_4 is connected (by parametric interaction at formation of coherent correlated state) with the optical phonon mode in PdD with the frequency $\omega_0^{(2)} = \Omega_4/2 = \Omega_2$!

The different situation takes place if there is a full-scale (maximally possible) change of the oscillator frequency $\omega_0(t)$ (18). In this case the process of formation of totally correlated state is possible at various actions on the system (including the use of low frequency $\Omega \ll \omega_0$). Figure 4 presents the time dependencies of the correlation coefficient $r(t)$ for full-scale (18) change in the oscillator frequency $\omega(t)$ at $\Omega = \omega_0/100\pi \approx 0.03\omega_0$.

Obtained results can explain Rossi–Focardi experiments at action of radio-frequency irradiation to hot NiH nano-powder situated in closed chamber with the presence of compressed H_2 gas [11,12]. In this case the action of RF-irradiation on surface of nano-particles leads to modulation of acoustic phonon and plasmon modes of these nano-particles. At such modulation the processes of formation of correlated states and stimulation of nuclear reactions

$$\mathrm{Ni}^A + \mathrm{p} \to \mathrm{Cu}^{A+!} + \nu$$

are possible. Barrier transparency for these reactions at temperature $T \approx 400 - 600$°C increases from $D_{r=0} \approx 10^{-1000}$ (for non-correlated states of interacting p and Ni^A nuclei) to $D_{r=0.999993} \approx 10^{-6} - 10^{-4}$ (for correlated state of p)!

The similar effect of formation of coherent correlated states of D^+ions and stimulation of effective nuclear dd-fusion

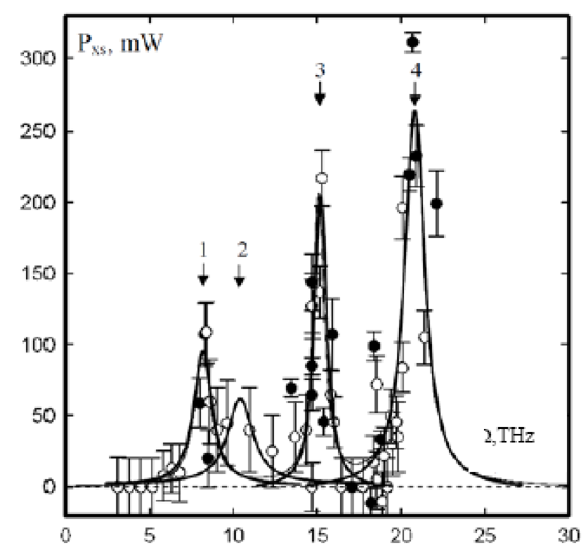

Figure 3. Frequency dependencies of energy release at combined action of beat frequency of two different lasers on surface of cathode during the electrolysis [9].

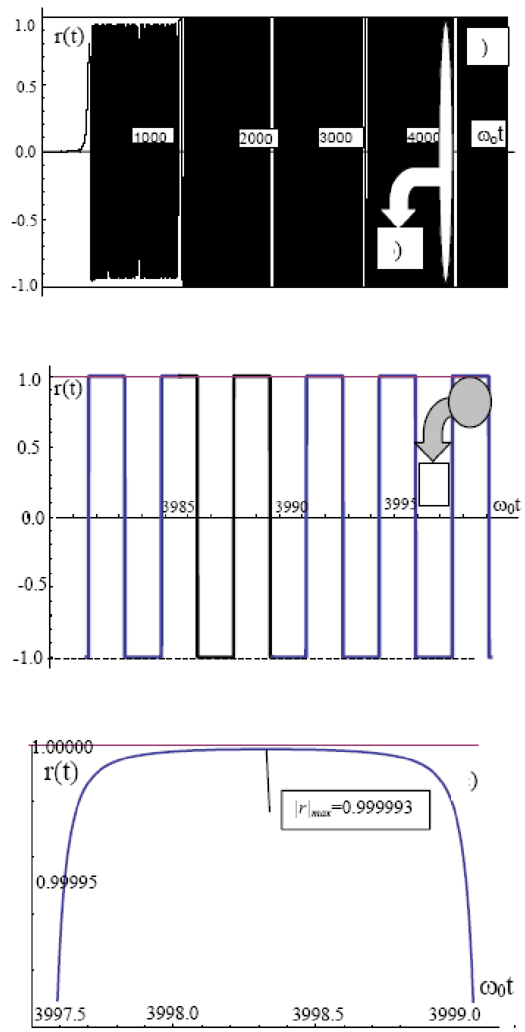

Figure 4. Dynamics of change of correlation coefficient during the action of parametrical perturbation with the frequency $\Omega = \omega_0/100\pi$ at full modulation $\omega(t) = \omega_0 \cos \Omega t$ of oscillator. Figure 4(b) corresponds to the increased fragment of Fig 4(a) and Fig. 4(c) corresponds to the increased fragment of the Fig. 4(b).

(including generation of neutron bursts) takes place in cooled D_2 gas at changing of strong external magnetic field in interval 8–10 Oe [16].

4. Summary

In conclusion, note that a coherent correlated state whose influence on the penetration probability into the region under the potential barrier is associated with the synchronization of the momentum fluctuations at different levels of the superposition state and the particle kinetic energy fluctuation caused by this synchronization can be realized not only in a parabolic field [2]. Such state can be formed, at least in principle, within any system of levels of quantized motion that is not subjected to an external intense dephasing action, provided that a certain coherent action is superimposed on it.

These effects may play the important role in processes of controlled nuclesynthesis in different potential wells and systems with nonstationary action on active systems (e.g. in experiments with high-current electron driver in Kiev Laboratory "Proton-21" [13] and in experiments on nuclear transmutation of stable and radioactive isotopes in growing biological systems [14,15]).

Another method of nuclear reaction optimization is connected with the process of creation of the spatially-compressed states of particle beam under action of crystal fields.

In ordered crystal lattice there is very strong influence of crystal axes and planes electrical field on motion and interaction of fast charged particles with crystal atoms and nuclei. It was shown in our old articles [17,18] that in monocrystal targets like LiD the rate of fusion process with the participation of both target nuclei (e.g. D) and beam of fast nuclei (e.g.T), directed at *Lindhard* angle, may be increased by 10–100 times compared to the alternative process of deceleration on atomic electrons. Such changes are based on the use of specific channeling physics regime of motion – "overbarrier motion" At such regime the processes of spatial redistribution and dechanneling of accelerated ions take place.

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